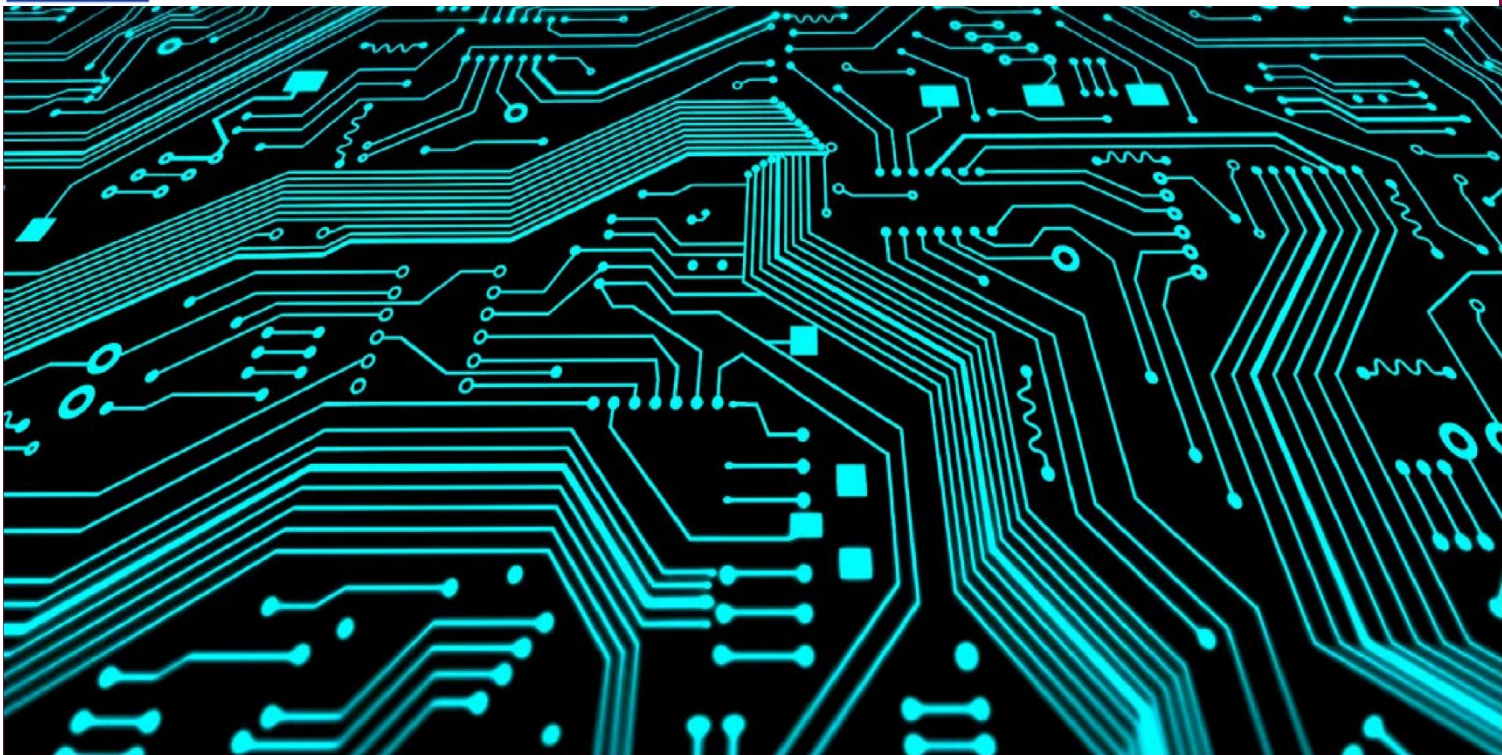




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Republic of Uzbekistan**

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OPTICAL AND PHYSICAL PROPERTIES OF SEMICONDUCTOR MATERIALS

Tashkent – 2025

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**OPTICAL AND PHYSICAL PROPERTIES OF SEMICONDUCTOR
MATERIALS. Study guide.**

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The textbook is based on a new program based on the analysis of international curricula and programs for higher education. It analyzes the properties of real materials, optical and electrophysical parameters while at the same time explaining the fundamental foundations of semiconductor physics. The general physical and optical parameters of various semiconductor materials are presented in multiple tables with the possibility of their mutual comparison.

The manual was prepared as one of the practical results of the Project "Development of the training of target engineers in solar energy in Uzbekistan" (101128871 - DEBSEUz - ERASMUS-EDU-2023-CBHE) within the framework of the European Union's "Erasmus+" program.

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INTRODUCTION

Semiconductor materials form the basis of the most common sources of photoelectric energy. In the process of using any energy sources, controlling their photoelectric and energy parameters, if necessary, correction as well as connecting them to electrical networks, electronic devices created on the basis of semiconductor materials are used. Therefore, trained specialists need to know electronic devices created on the basis of semiconductor materials, their operating specifications, operating parameters, and the procedure for connection to electrical circuits. Before acquiring such knowledge, students are required to familiarize themselves with the optical and physical properties of semiconductor materials.

Today, the types of electronic devices created on the basis of semiconductor materials are very wide, since the exponential characteristics of the electronic device and the physical parameters that represent them are limited by the properties of the material used as a basis. Among such devices, optoelectronic and photoelectronic devices, including the optical and electrophysical properties of materials from which solar cells are made, are considered extremely important. When thinking about the properties of optics, it is important, first of all the width of the material's forbidden energy zone, whether the optical transitions are correct or not. The crystallographic quality present in the material volume, the presence of various defects, their possible negative or positive influence on charge migration processes when light is incident and temperature change can cause variation of optical and physical properties. In addition, the physical and optical properties of thin plates or thin layers obtained from the material for which the optical or photoelectric device is prepared are related to their surface state. The application of technologies, such as using other types of optical materials to reduce recombination sensitivity by improving surface conditions or the formation of textures of different shapes and sizes on the surface can lead to improvement of properties, energy parameters of semiconductor-based optical and electronic devices.

The versatility of semiconductor materials, aside from elementary semiconductor materials such as silicon and germanium, show competition in solar cells created on alternative materials such as atomic alloys consisting of various atomic compounds, molecular alloys based on the deposition of molecules, organic and polymer materials.

Thanks to the use of modern nanotechnology achievements, the possibility of radically improving the output parameters of solar cells, including electronic devices created on the basis of semiconductor materials, appears. By applying quantum-sized effects such as nanoplasmonic effect, boiling charge carrier effect to photovoltaics, radical improvement in optical and photoelectric properties of

semiconductor materials and ultimately solar cells created on their basis is achieved. In connection with this, a new training course "Optical and physical properties of semiconductor materials" was formed. This Study Guide is based on a new curriculum developed by analyzing international curricula and programs. It analyzes the properties of real materials, optical and electrophysical parameters, while at the same time explaining the fundamental foundations of semiconductor physics. The general physical and optical parameters of various semiconductor materials are presented in multiple tables with the possibility of their mutual comparison.

The manual is presented as one of the practical results of the project "Development of the training curriculum for target engineers in solar energy in Uzbekistan" (101128871 - DEBSEUz - ERASMUS-EDU-2023-CBHE) within the framework of the European Union's Erasmus+ program. The manual has 14 chapters and is 424 pages long. At the end of each chapter, there are questions for the control, and at the end of the manual, there are test samples and a "Glossarium." It is advisable to use the manual for teaching courses in the areas of Bachelor's "Alternative energy sources", "Physics", "Physics of semiconductors", "Physics of solids", "Physics of condensed matter", etc.

Chapter I

SEMICONDUCTORS.

ELEMENTARY THEORY OF ELECTRIC CONDUCTIVITY

1.1. Classification of substances according to their specific electrical conductivity. Semiconductors

All substances can be divided into three broad categories according to their electrophysical properties: metals, semiconductors, and dielectrics. Substances are generally classified according to their specific electrical resistance (ρ). The specific electrical resistance of metals is in the range of 10^{-6} - 10^{-4} $\Omega \cdot \text{cm}$ (e.g., $\rho = 1.58 \cdot 10^{-6}$ $\Omega \cdot \text{cm}$ for gold, $\rho = 1.05 \cdot 10^{-4}$ $\Omega \cdot \text{cm}$ for nickel alloy). Substances with an electrical resistance of 10^{-4} to 10^{10} $\Omega \cdot \text{cm}$ are called semiconductors. Finally, substances with an electric resistance greater than 10^{10} $\Omega \cdot \text{cm}$ are called dielectrics (e.g., at room temperature, the electric resistance for mica is 10^{13} - 10^{16} $\Omega \cdot \text{cm}$, and for glass 10^8 - 10^{15} $\Omega \cdot \text{cm}$).

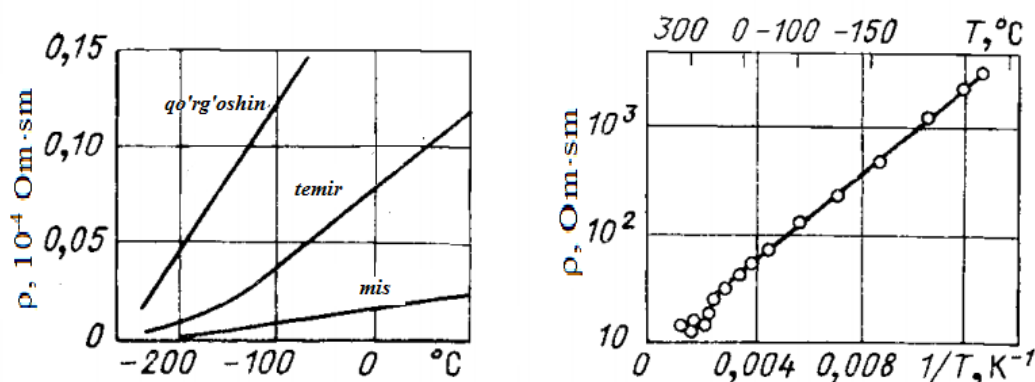


Fig. 1.1. The temperature dependence of the specific resistance of metals and silicon

From the given examples, it can be seen that the values of electric resistance in the transition from one class to another do not exactly coincide. Therefore, electric resistance values cannot be an exact criterion for the classification of substances. As can be seen from Fig. 1.1 and 1.2, metals and semiconductors have different temperature dependencies of electrical resistance and electrical conductivity. For metals, with increasing temperature, ρ increases proportionally, i.e.

$$\rho = \rho_0(1 + \alpha t) = \rho_0 T / T_0, \quad (1.1)$$

where, ρ is the specific resistance of the metal at $\rho_0 - 0^\circ\text{C}$; α - is the coefficient of resistance by temperature, which is equal to $1/273$; $T_0 = 273$ K.

In metals, according to Fig.1.2, the dependence of specific electrical conductivity on temperature is negative, which means that with an increase in temperature, its conductivity decreases.

For semiconductors, the temperature dependence of the specific electrical resistance and conductivity is different, and over a given temperature range they are represented by the following equations:

$$\rho = \rho_0 e^{\beta/T} \quad , \quad (1.2)$$

$$\sigma = \sigma_0 e^{-\beta/T} \quad , \quad (1.3)$$

where ρ_0 , σ_0 , β - are certain constants that are specific to each semiconductor. Such dependencies are characteristic of semiconductors, which are unpredictable with respect to the temperature of electrical resistance and conductivity. Then, according to Figure 1.2, the dependence of specific conductivity on temperature is proportional, i.e.

$$\frac{\Delta\sigma}{\Delta T} = \frac{\sigma_2 - \sigma_1}{T_2 - T_1} > 0 \quad (1.4)$$

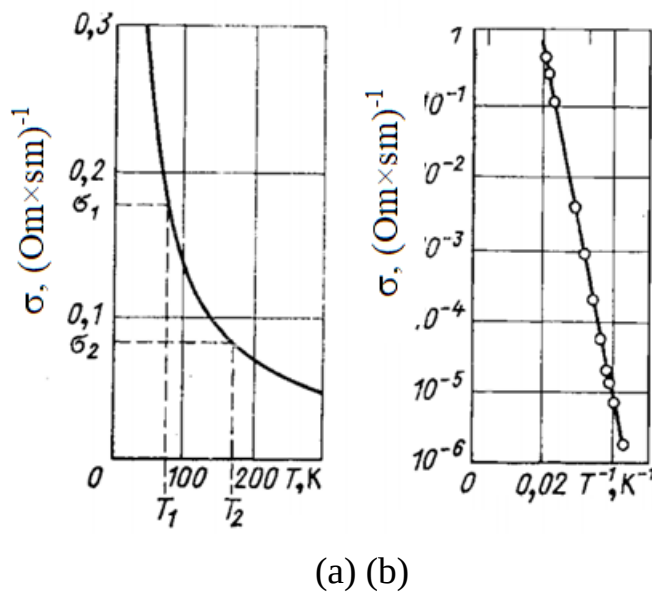


Fig. 1.2. Temperature dependence of lead and silicon specific conductivity

It can now be thought that the difference between semiconductors and metals can be solved by the temperature-dependent conductivity coefficient designation. However, in a certain range of temperature, a semiconductor can show a similar bond to a metal. Therefore, it is not always possible to determine whether a substance falls into the class of semiconductors by the sign of the dependence of specific conductivity on temperature.

The answer to this question can be given if we observe how its permeability changes during the process of cooling of a substance. According to Fig.1.2, these variations can vary.

As the temperature decreases, the specific conductivity of metals increases, this is possible only because metals have always have free electrons present depending on the temperature. For semiconductors, on the other hand, the specific conductivity decreases as the temperature decreases (Fig. 1.2-b). This means that the appearance of free charge carriers in a semiconductor occurs solely due to the effect of heat. These charge carriers are called heat or balanced charge carriers. Experiments show that free charge carriers in a semiconductor appear not only under the action of heat, but also in light, radiation with nuclear particles. The charge carriers formed in these cases are called unbalanced charge carriers. The formation process of balanced and unbalanced charge carriers depends on the structure of the semiconductor material and the presence of compounds in its composition.

Thus, semiconductors have a specific conductivity of 10^{-10} to 10^4 $(\Omega \cdot \text{cm})^{-1}$ at room temperature, which strongly depends on the structure of matter, the type and quantity of the mixture, and external conditions: temperature, light, radiation from nuclear particles, the influence of electric and magnetic fields.

As you know, there are 2 types of conductors of electric current: electron and ionic. Metals - are conductors with an electron. In ionic conductors, current passes through with the help of ions, so when an electric current passes through an ionic conductor, its composition changes. Typical representatives of ionic conductors are - electrolytes.

Semiconductors can be both electron and ionic. This book will only deal with the physical properties of electron semiconductors.

Electron semiconductors contain many different substances. Simple substances as semiconductors: boron *B*, silicon *Si*, phosphorus *P*, sulfur *S*, germanium *Ge*, arsenic *As*, gray tin *Sn*, surfacing *Sb*, selenium *Se*, tellurium *Te* It also contains many complex chemical compounds. As an example, Table 1.1 shows the combinations of some elements of Mendeleev's periodic table with semiconductor properties (Figure 1.3). The table shows the crystal structure, lattice constant, and other properties of these important semiconductors.

The electrophysical properties of semiconductors have enabled their wide application in the creation of various devices. Semiconductor devices are used to convert alternating current to direct current (diodes), change or move frequencies, amplify and generate electrical oscillations (transistors, SVM diodes, Gann diodes, lavina-conductive diodes, tunneling diodes, etc.).

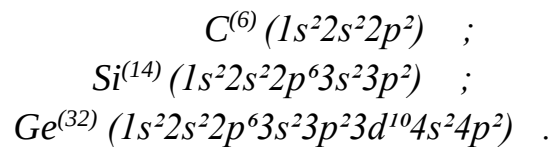
There are semiconductor devices such as semiconductor diodes, used to convert heat energy and radiation energy into electricity (thermoelectric elements and photocells), to convert electrical energy into light (phosphor, light-emitting diodes, lasers). Semiconductor devices are also used as "sensors", e.g. for measuring temperature (thermistors), for detecting light and particle radiation (photoresistors

and dosimeters), for measuring pressure (transducers) and for measuring magnetic fields (Hall sensors).

At the heart of the operation of each device lies a certain physical process and phenomena. In other words, semiconductor physics constitutes the theoretical basis of semiconductor electronics. Development of semiconductor materials, creation of new types of devices on their basis, conducting research in this field, production of high-quality semiconductor devices and integrated circuits, creation of new devices suitable for various applied conditions - all this requires in-depth knowledge of semiconductor physics.

1.2. Model representations about the conductivity mechanism of a semiconductor

Let's take a look at the electrical conductivity mechanism of uniatomic semiconductors, e.g., silicon and germanium. Silicon and germanium are in the same group as carbon in the periodic table of Mendeleev. Their circuits are distributed as follows:



The outer electron shell of these atoms is partially filled, each containing 4 electrons: 2 in the *s*- and 2 in the *p*- states.

In crystal formation, e.g. in silicon, four valent electrons of each atom pass from the $3s^2 3p^2$ states to the hybrid $3sp^3$ state, where anti-parallel spins are formed and form four covalent bonds. As a result, each atom is surrounded by four close neighbors, which are located in the center of the tetrahedron. Such a crystalline structure is cubic and is called a "diamond-type" crystal lattice (Fig. 1.4). In this structure, the four outer electrons of each atom form covalent bonds (two electrons for each bond). These bonds are depicted in Fig. 1.4 as directions connecting two adjacent atoms. Fig. 1.5, Fig. A, shows a two-dimensional image of covalent bonds in silicon crystal structure. Here, in the lattice node, the silicon ions have a +4 charge and have four valent electrons. The valent electrons that provide the covalent bond are shown by black dots in Fig. 1.5, a.

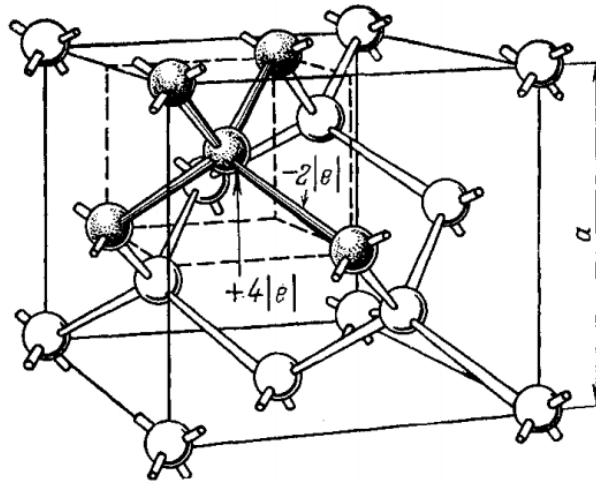


Fig.1.4. Diamond-type crystal lattice

In a semiconductor with an ideal structure, all electrons will be in the bonded state, as shown in Fig. 1.5. If such a semiconductor is placed in an external electric field, no electric current will be generated, since all covalent bonds in the grid are completed and there are no free charge carriers.

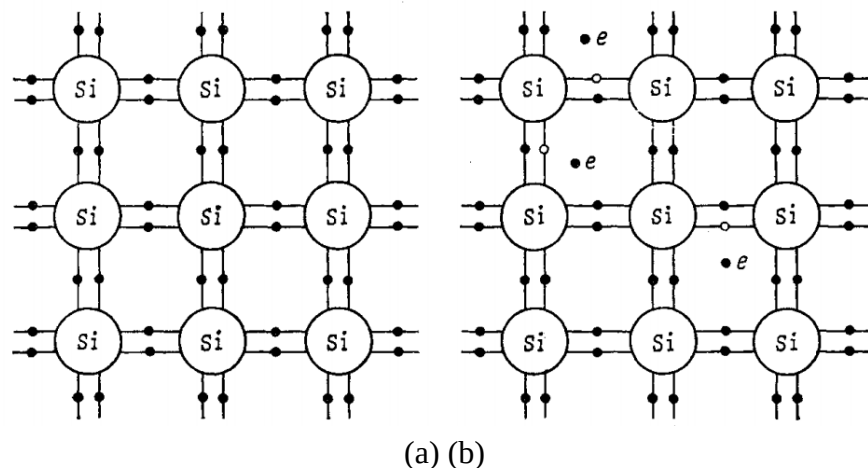


Fig. 1.5. Two-dimensional view of the gardens on the silicon lattice

Suppose that if some external force, such as heat, breaks the covalent bond in a semiconductor and transfers the electron into a free state. This process of transferring a bonded circuit into a free circuit state is called "generation." When the electron covalent leaves the bond, the bond is incomplete. There will be an excess positive charge of the false bond because it is no longer compensated by the charge of the electron. The space in the covalent bond is called the "hives." In general, the crystal electron remains neutral because each generated positive charge is filled with a corresponding free electron. In Fig.1.5-b, the free electrons and sleds are depicted as black and bright circles, respectively. A free electron can occupy space in the covalent bonding and return to the bonded state. This process converts a free electron into a bonded circuit state called "recombination."

Thus, a pure semiconductor - is an ideal pure substance in which an equal number of free electrons and holes are formed as a result of the violation of covalent bonds. At room temperature, the concentration of free electrons and holes in germanium is about 10^{13} cm^{-3} , and in silicon it is about 10^{10} cm^{-3} . Increasing the temperature of the semiconductor leads to an increase in the number of covalent bonds, and at the same time, an increase in the concentration of free electrons and holes in the semiconductor.

In the absence of an external electric field, a semiconductor crystal carries out a free electron thermal movement, collides with the defects in the crystal lattice, as a result of which its direction of motion changes. Therefore, the thermal motion of a free electron is random, as illustrated in Fig. 1.6-a.

The space in the covalent bond can be filled by the electron passage from the adjacent covalent bond as a result of thermal excitation. In the process of such crossing, the hole also moves randomly.

Thus, the thermal motion of free charge carriers in a semiconductor crystal is random. Each charge carrier illustrates a complex direction of motion (Fig. 1.6-a).

The distance between two free-charge carrier collisions is called the "free pass length," and the average value of all motion segments is the "average free pass length." Accordingly, the time between two collisions and its mean value is referred to as "free run time" and "mean free run time".

The mean free pass length l and the mean free run time τ are related by the relation $l = v_0 \tau$ below, where v_0 - is the average velocity of the thermal motion of the free charge carrier (the average value of the velocity modulus). In semiconductors, the velocity of thermal movement of electrons at room temperature is approximately 10^7 cm/s .

Because the thermal motion of free charge carriers - electrons and holes - are random, their average velocity as a vector is zero. This means that for every moving free electron (holes), it is possible to find a free electron (holes) moving in the opposite direction at the same speed (Fig. 1.7-a). Thus, as a result of random thermal motion, the number of free electrons and holes in the direction in which the holes are moving will be equal to the number of electrons and holes moving in the opposite direction, on average. Therefore, in the absence of an external electric field, the total charge propelled by free electrons and holes becomes zero, and the random thermal motion does not bring current to the charge carriers.

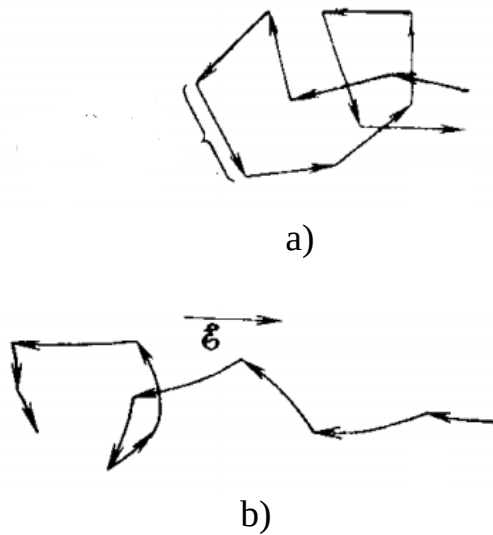


Fig. 1.6. In the absence of an electric field a) and in the presence of b) the motion of a free electron

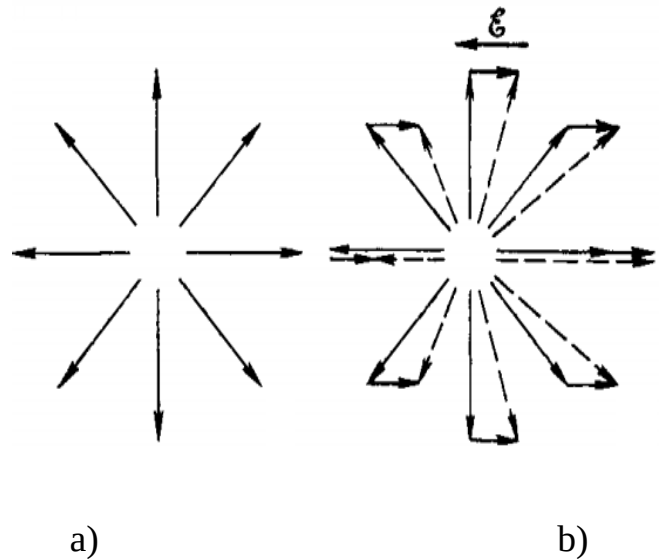


Fig. 1.7. In the absence of an external electric field, a) and b) the velocity of conduction electrons in the presence of

If a particular semiconductor is in an external electric field, then under the action of this field, the free electrons acquire a velocity directed towards the field. By doing so, as a result of thermal movement, the velocity of electrons moving against the external electric field increases, and the velocity of electrons moving in the direction of the field decreases.

As a result, the entire sum of the free electrons moves at a certain speed against the external electric field (Fig. 1.7-b). The change in the speed of motion of electrons affects their kinetic energy. The electrons colliding with the defects of the crystal lattice transfer all the energy received in the field to the lattice, and their energy comes into thermal equilibrium with the lattice. After the collision, the electrons, continuing their random heat motion, again acquire a directed motion against the external electric field.

Thus, the actual motion of an electron in a crystal consists of an orderly motion caused by the action of a random heat motion and an external electric field (Fig. 1.6-b). As a result, the entire sum of the free electrons travels at a given average speed. The directed motion of charge carriers in the outer electric field is called "drift," and their directed motion velocity is called "drift velocity." Fig. 1.7-b depicts the vectors of the rate of drift with horizontal lines, and the vectors of thermal motion and the total velocity of drift with dot lines. In this case, the average speed of motion of the free electrons is not zero, and a current can flow in the direction of the electric field in the semiconductor, which is induced by free electrons.

The electrons in the filled bonds, under the influence of an external electric field, move against the field direction when they pass into a space in a covalent bond. At the same time, the space in the covalent bond is moving in the direction of the

external electric field, which is equivalent to the movement of a positive charge across the field. Such an electrophysical process, i.e., the movement of electrons to the empty spaces in the covalent bond and the movement of holes across the outer field, is called "hole conductivity."

Thus, in a pure semiconductor, i.e., in the absence of mixtures, electron and hole conductivity takes place. This means that in a pure semiconductor, the electric current is determined by the flow of electron and hole currents from two components in the same direction.

The permeability of pure semiconductors can be explained from an energy point of view. In isolated atoms, electrons have, according to quantum theory, discrete energy values. At each energetic level, there can be only two electrons according to the Pauli principle, and they must have opposite spins. When semiconductor materials are formed, that is, when atoms approach a distance of about 10^{-8} cm, the wave functions of atomic electrons cancel each other out.

Therefore, the energy surface of the valent electrons passes through the zone form. This zone is called the "valent zone". From the permeability of atoms at a higher level, a new zone is formed, which is called the "conduction zone." Each zone occupies a certain energy range. Fixed energy zones are separated from each other by an interval called an "energy gap" or "forbidden zone". Fig.1.8 shows the energy diagram of a special semiconductor, where E_c - denotes the minimum energy of the free electrons in the crystal, E_v -- the maximum energy of the electrons in the valent band, E_g the width of the forbidden band, this value depends on the type of chemical bonding and the atoms that form the crystal lattice.

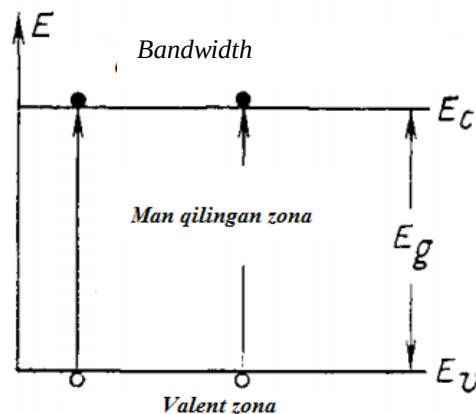


Fig. 1.8. Energy Zones of Pure Semiconductor

As you know, current density - is the charge that is transported by electrons through a unit surface in a unit of time. Therefore, taking into account the motion of one electron, the current density can be expressed as follows:

$$J = -\frac{e}{\Omega} v, \quad (1.5)$$

where v - is the average velocity of the electron, Ω -volume.

If all possible quantum states are filled with electrons in the energy zone, then the density of the electric current in the system

$$J = -\frac{e}{\Omega} \sum v_s + \frac{e}{\Omega} v_i \quad (1.6)$$

Across the entire zone, an electron with velocity v for each electron can be filled with an electron with velocity $-v$ in the opposite direction, which forms a current that is equal but the direction is reversed.

$$J = \frac{e}{\Omega} v_i \quad (1.7)$$

A breakdown of the covalent bond in a private semiconductor, for example, leads to the formation of a gap in the free electron and valent bond under the influence of heat. According to the energy zone diagram shown in Figure 1.8, this is equivalent to the passage of an electron from the valent band towards the permeable zone. In this case, all the quantum states in the valent band would be filled with electrons, except for one state. The motion of the charge carrier, in the transition to the free state in the covalent bond, is determined by v_s . Then the current density of all electrons in the valent band is expressed as

$$J = -env \quad (1.8)$$

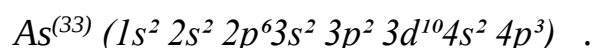
However, the sum of the current densities in the 1.7 expression is zero across the zone. With this, the sum total of all electrons in the valent band is equal if we place the last electron in the space in the covalent bond and add a positive charge $+e$ to it. Such an empty position is called a "positive stake," or simply a "spit."

At normal temperatures, in the conductive zone of a particular semiconductor, there will always be some electron emitted by heat from the valent band, and in the valent band there will be an equal number of holes with them. If there is an external electric field, then free electrons and holes produce an electric charge conduction.

1.3. Model representations about the mechanism of electric conductivity in semiconductors with input

Now, one of the atoms in a silicon crystal belongs to the V group, and we look at the mechanism of electrical conductivity in a semiconductor replaced with arsenic, for example. A semiconductor that has mixtures is called an inlet conductivity, and its electrical conductivity is due to the presence of the mixture in the crystal, which is called inlet conductivity.

The arsenic atom has five valent electrons, which are located in the $4s$ - and $4p$ states:



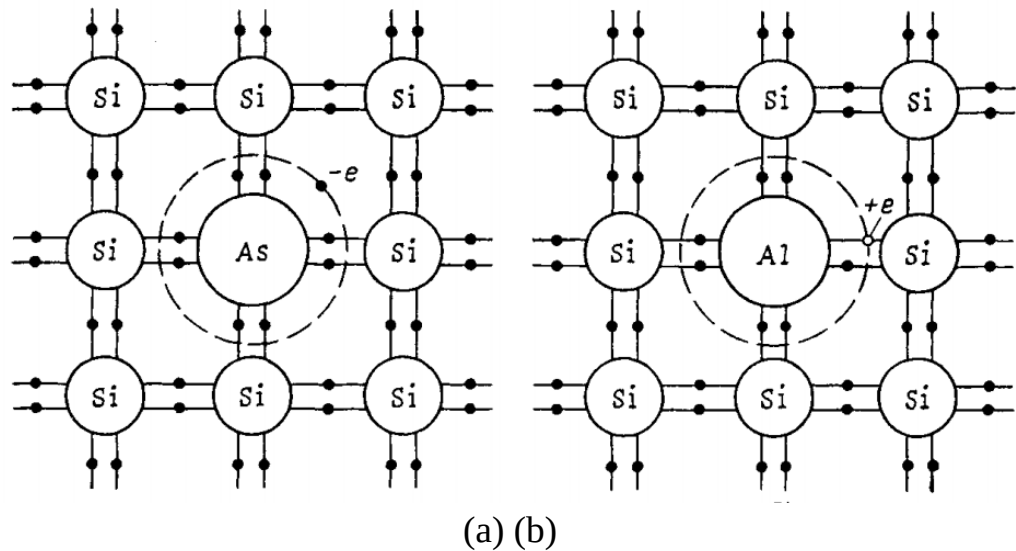


Fig. 1.9. Two-dimensional view of bonds in a donor and acceptor semiconductor

In silicon crystal, the four valent electrons of arsenic form a covalent bond together with the four valent electrons of silicon atoms, which is shown in Fig. 1.9-a. The fifth electron of arsenic, uncoupled to a covalent bond, is weakly bonded, where the electrons around the silicon atom settle in the electric field.

At low temperatures, the fifth electron is localized around the arsenic atom, but at high temperatures, this electron can detach from the mixture and move freely throughout the crystal. In this case, too, the electron neutrality of the crystal is preserved, since the arsenic atom, giving its fifth electron, turns into a positive ion.

Along with ionization of the mixture, the atoms of the underlying substance can be ionized. However, at temperatures below the temperature required for a significant increase in electrical conductivity, the number of electrons separated from the mixture becomes much greater than the number of electrons and holes formed by the breaking of covalent bonds. Therefore, electrons play a dominant role in the conductivity of a semiconductor, which is why they are called primary charge carriers, while holes are called non-fundamental charge carriers. Such a semiconductor is called an electron or n-type, and the mixture that gives electrons is called a donor or n-type compound.

The presence of a mixture in a semiconductor crystal leads to the emergence of local energy layers in the crystal, which settle within the negative zone. The ionization of the arsenic atom results in the formation of a free electron, and much less energy is required to create it than breaking the covalent bond, so that the energy level of the donor mixture is located very close to the boundary of the permeable zone within the negative zone (Fig. 1.10-a).

Suppose that atoms of a third group are added to the covalent bonding lattice of a semiconductor crystal, for example, aluminum atoms are added to a silicon

lattice. The high stress of aluminum is three: $Al^{13} (1s^2 2s^2 2p^6 3s^2 3p^1)$, so one bonding of a silicon atom would not be complete (Fig. 1.9-b).

As a result of thermal action on the incomplete bonding around the aluminum atom, an electron can pass from the silicon atom to the aluminum atom. At the same time, a negative charge of the aluminum ion and a free hollow is formed, which is involved in the conductivity of the semiconductor by moving in silicon bonds. Mixtures that have electron attraction properties are called acceptor mixtures. The energy required for the electron to pass from the silicon atom into the mixture of the receiver is much less than the energy required to break the covalent bond, so the number of hives will be much greater than the number of free electrons, and therefore in such a semiconductor, the main charge carriers will be hives, and the electrons will be non-basic charge carriers. A semiconductor with an acceptor mixture is called a "hollow" or "p-type."

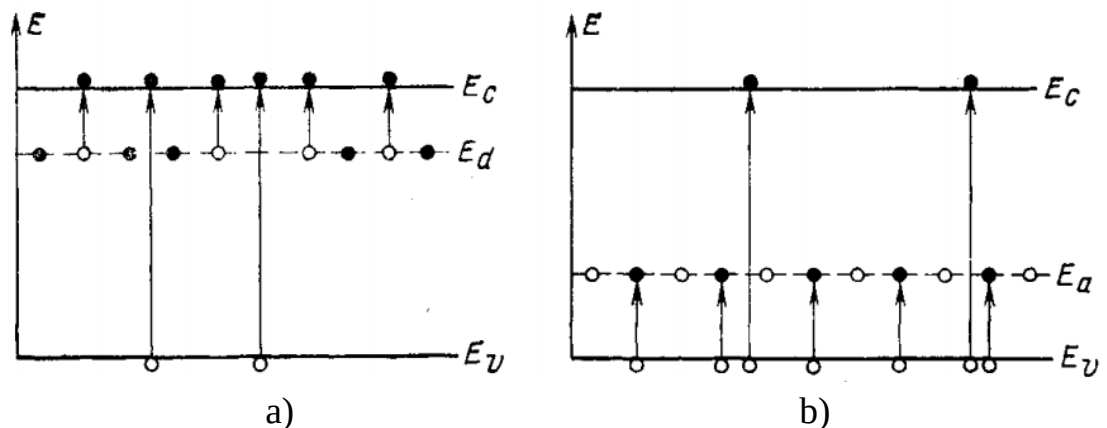


Fig. 1.10. Energy diagram of donor and capacitor semiconductor

According to the energy diagram shown in Fig. 1.10 b, the energy level of the acceptor mixture is located slightly above the upper limit of the negative zone E_a . As a result of ionization of the acceptor mixture, the electron passes from the valent band to the E_a state, and a space - cavity is formed in the covalent bond in the valent band, which is a free charge carrier.

In semiconductors, mixtures of donor and acceptor may exist simultaneously.

1.4. Elementary theory of electrical conductivity in semiconductors

Since we perform a current density calculation in a donor semiconductor, we consider electrons as ideal particles, they do not have their own size, and do not interact with each other. According to the classical theory, the radius of an electron r_0 is about 10^{-13} cm, so if their concentration is $n \sim 10^{22} \text{ cm}^{-3}$, the size of the

electrons will be $r \sim 1/n$, and the first assumption is considered correct. Quantum theory of a solid, as shown in Chapter 2, allows us to treat electrons in a crystal as non-interacting particles.

Suppose that the concentration of electrons n (the number of free electrons per 1 cm^3 part of a semiconductor) and their drift velocity is v . Current density is defined by electrons as the passage of a charge through a unit field in a unit of time, therefore:

$$J = -env \quad (1.9)$$

Now we determine the drift velocity of the electrons. Suppose that dt/τ is the probability τ that an electron will collide in time dt . The probability of a collision does not change in terms of time $1/\tau$, i.e. τ - is a constant quantity.

At the time dt where the number of electrons is n , the number of collisions is equal to ndt/τ . At the same time, the number of charge carriers in time dt decreases by n as follows:

$$-dn = n \cdot dt/\tau \quad (1.10)$$

Solving this differential equation, the collision-free motion of electrons can be expressed as follows:

$$n(t) = n_0 e^{-t/\tau}, \quad (1.11)$$

where n_0 is the number of electrons in $t=0$.

The external electric field E gives acceleration to electrons

$$a = \frac{eE}{m} \quad (1.12)$$

Electronic during free run time

$$v_t = at = \frac{etE}{m} \quad (1.13)$$

The drift will have velocity, and

$$x = \frac{eE}{2m} t^2 \quad (1.14)$$

goes the distance.

The distance that all electrons travel in the direction in which they correspond to the external electric field will be equal to

$$X = \frac{en_0 \tau^2 E}{2m} \int_0^\infty \left(\frac{t}{\tau}\right)^2 e^{-t/\tau} \frac{dt}{\tau} = \frac{en_0 \tau^2 E}{m} \quad (1.15)$$

If electrons τ have an average free travel time, then

$$T = n_0 \tau \quad (1.16)$$

Drift velocity of electrons is defined by expression.

$$v = \frac{X}{T} = \frac{e \tau E}{m} \quad (1.17)$$

This means that the speed of drift is related to the time of movement of the electrons, which is proportional to the voltage of the external electric field, the time of free travel, and inversely proportional to the mass of the electron.

The parameter that relates the rate of drift of charge carriers to the voltage of the external electric field is called "charge carrier mobility". We define it with μ , so:

$$v = \mu E, \quad (1.18)$$

$$\mu = v/E = e\tau/m \quad (1.19)$$

where μ is the mobility of the charge carriers, that is, the drift velocity of the charge carriers at one unit of voltage.

Taking into account the formula (1.19), the expression for the density of (1.9) can be written as follows:

$$J = -env = en\mu E, \quad (1.20)$$

Here the velocity of the electrons moves in the opposite direction to the voltage vector v .

Using the expression (1.20) and Ohm law, we can write the mean conductivity as follows:

$$\sigma = J/E = en\mu \quad (1.21)$$

specific conductivity will take on a look.

$$\sigma = e^2 n \tau / m \quad (1.22)$$

If we subtract τ from equation (1.4), the relation (1.22) can be written as:

$$\sigma = e^2 n l / m v_0 \quad (1.23)$$

However, it is really necessary to take into account the total velocity of the electron when determining τ , which is determined as the sum of the thermal velocity v_0 and the rate of drift v_d , i.e.:

$$v = v_0 + v_d, \quad (1.24)$$

$$\tau = l/v. \quad (1.25)$$

Because the speed of the electron's drift depends on the strength of the external electric field, its full velocity $v = v_0 + v_d$ and therefore the free travel time τ depends on the energy of the particle. Therefore, to determine the drift velocity, it is necessary to average the free run time by taking into account the distribution function of the electrons in terms of energy. This issue will be dealt with in Chapter 6.

Control Questions

1. What is the role of semiconductors in modern physics and technology?
2. What are semiconductors?
3. How are substances classified according to their specific electrical conductivity?

4. What happens to the electrical conductivity of p-type and n-type semiconductors?
5. What is a donor impurity?
6. What is an acceptor impurity?
7. What does the electrical conductivity of semiconductors depend on?

Chapter II

FUNDAMENTALS OF SEMICONDUCTOR ZONE THEORY

2.1. Schredinger's equation for crystal

As you know, solids are made up of atoms, atoms are made up of nuclei, and electrons. The nuclei of atoms form a crystal lattice, which has the property of spatial periodicity. Although atomic nuclei are charged, the lattice is practically not deformed by the action of the external electric field. This means that the forces that hold atomic nuclei at the nodes of the crystal lattice are usually much stronger than the forces generated by external electric fields.

Electrons that are not located close to the atom and therefore are not strongly attracted to it move through the crystal, generating an electric current. However, serious difficulties arise in the quantitative description of this phenomenon. The difficulty arises from the fact that electrons are charged particles and collide with other electrons in their motion. But there are electron repulsion forces between electrons, which connects the motion of the electron with the motion of the electrons around it. That is, in this case, the problem of not one electron, but many electrons must be solved. Therefore, in order to determine the stationary states and energy spectrum of many atomic nuclei and electron systems in a crystal, it is necessary to solve the Schrödinger equation:

$$H\Phi = E\Phi, \quad (2.1)$$

Here H is the Hamilton operator of the crystal; Φ - is its wave function; E - crystal energy.

The wave function of a crystal depends on the coordinates r_i of all electrons and R_l of all atomic nuclei:

$$\Phi = \Phi(r_1, r_2, \dots, r_n, R_1, R_2, \dots, R_N).$$

Hamilton's operator includes:

1. The operator representing the kinetic energy of electrons is, $\sum_{\alpha} \left(-\frac{\hbar^2}{2m_0} \Delta_i \right)$, where $\hbar$ is the Planck constant, m_0 - electron mass, Δ_i - Laplace's operator;
2. The operator representing the kinetic energy of nuclei is, $\sum_{\alpha} \left(-\frac{\hbar^2}{2M_{\alpha}} \Delta_{\alpha} \right)$
3. The potential energy of the paired interaction between electrons-
$$\frac{1}{2} \sum_i \sum_j \frac{e^2}{r_{ij}}$$
4. The potential energy of the paired interaction between nuclei-
$$V_0(R_1, R_2, \dots, R_N)$$

5. The potential energy of the interaction between electrons and nuclei-
 $U = V_0(R_1, R_2, \dots, R_N)$

Taking into account these components, the Schrödinger equation is expressed as follows:

$$\left\{ \sum_{\alpha} \left(-\frac{\hbar^2}{2m_0} \Delta_i \right) + \sum_{\alpha} \left(-\frac{\hbar^2}{2M_{\alpha}} \Delta_{\alpha} \right) + \frac{1}{2} \sum_i \sum_j \frac{e^2}{r_{ij}} + V_0(R_1, R_2, \dots, R_N) + U(r_1, r_2, \dots, r_n, R_1, R_2, \dots, R_N) \right\} \Phi = E \Phi \quad (2.2)$$

The number of independent variables in the equation depends on the total number of particles in the crystal, and in case of 1 cm³ substance they are approximately 10²³. Such a problem is not currently solved in a general form, but with the help of several successive approximations, its approximate solution can be obtained.

2.2. Adiabatic convergence and valent convergence

We divide the entire particle system into light (electrons) and heavy (atomic nuclei) particles. In the equilibrium state, the average values of the kinetic energy of these particles will be the same. The mass of the nucleus is much greater than the mass of the electron, i.e. $M_{\alpha} \gg m_e$, so the speed of motion of electrons is much higher than the velocity of the nucleus (about two orders of magnitude difference). When the state of any atomic nucleus changes, the spatial distribution of electrons adapts almost immediately to the new nuclear state. Therefore, the movement of electrons is independent of the motion of the nuclei, and there is no energy exchange between electrons and nuclei. This, at the first approximation, allows us to consider the motion of electrons in the potential field of strengthened nuclei. In this case, the wave function and energy of the electrons will be functions that change adiabatically with the change in the position of the nuclei, in which the coordinates of the nuclei participate as parameters in these functions.

If the motion of nuclei is studied, it is necessary not to consider the position of electrons, but to consider the field created by their mean spatial distribution. Such an approximation is called **the adiabatic** or **Born-Oppenheimer approximation**.

Let's use the simplest approximation, in which we assume that the nuclei are in a state of rest. In this case, the core coordinates $R_1, R_2, \dots, R_N$ are no longer variables, but the coordinates at the nodes of the crystal lattice become parameters as $R_{10}, R_{20}, \dots, R_{N0}$. When the nucleus is in a state of rest, $R_{\alpha} = R_{\alpha 0}$, the kinetic energy of the nuclei becomes zero, and the energy of their interaction V_0 takes a constant value. It can be zeroed by selecting an energy starting point. With this, equation (2.2) simplifies and now only describes the motion of electrons in nuclei in a resting state.

In this case, the energy of the electrons is E_e and their wave function Φ_e is now parametric dependent on the location of the nuclei $R_{\alpha 0}$.

As a result, the Schrödinger equation for electrons moving in at rest state nuclei is written as follows:

$$\left\{ \sum_{\alpha} \left(-\frac{\hbar^2}{2m_0} \Delta_i \right) + \frac{1}{2} \sum_i \sum_j \frac{e^2}{r_{ij}} + U(r_1, r_2, \dots, r_n, R_1, R_2, \dots, R_N) \right\} \Psi_e = E_e \Psi_e \quad (2.3)$$

But the number of variables in this equation is also so large that it cannot be deciphered.

2.3. One-electron approximation

One of the most common ways to solve a multi-electron problem for a crystal is the Hartree-Fock method, which allows you to turn a multi-electron problem into a single-electron problem. His idea is that the paired interaction between electrons is replaced by the average effect of each electron with all the other electrons.

Suppose i -th is the potential energy of electron Ω_i , which depends not only on its motion of all other electrons, but also on its own motion, since the motion of the electron affects the motion of the other electrons. Given the motion of electrons and the interaction between them, this field is termed as the "self-adjusting" field. By introducing the self-adjusting field, we convert the pair sum in equation (2.3) into only one sum, which includes the effect that each electron depends only on its own coordinates

$$\frac{1}{2} \sum_i \sum_j \frac{e^2}{r_{ij}} = \sum_i \Omega_i(r_i) \quad (2.4)$$

Similarly, the potential energy of the interaction between electrons and nuclei can also be expressed in aggregate form:

$$U(r_1, r_2, \dots) = \sum_i U_i(r_i) \quad (2.5)$$

With all of this in mind, we simplify equation (2.3) as follows:

$$\sum_{\alpha} \left(-\frac{\hbar^2}{2m_0} \Delta_i + \Omega_i + U_i \right) \Psi_e = E_e \Psi_e, \quad (2.6)$$

$$\sum_{\alpha} H_i \Psi_e = E_e \Psi_e, \quad (2.7)$$

Here H_i - is the Hamiltonian operator of the i -th electron, and its expression would be as follows:

$$H_i = -\frac{\hbar^2}{2m_0} \Delta_i + \Omega_i + U_i, \quad (2.8)$$

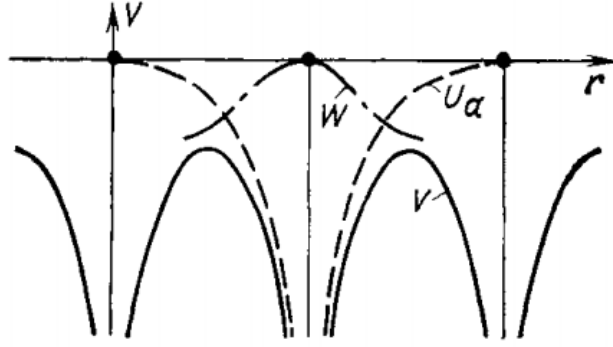


Fig. 2.1. One-dimensional grid

In this way, a multi-electron matter can be converted into a single-electron matter by means of an embedded self-adjusting field. By means of this approximation, electrons are treated as independent particles, meaning that they do not interact with each other.

$$\Psi_e(r_1, r_2, \dots) = \Psi_1(r_1) \Psi_2(r_2) \dots = \prod \Psi_i(r_i), \quad (2.13)$$

$$E_e = E_1 + E_2 + \dots = \sum E_i, \quad (2.14)$$

Schredinger's equation for each electron

$$H\Psi = E\Psi \quad (2.15)$$

looks. If

$$V(r) = U(r) + \Omega(r) \quad (2.16)$$

specify, the Schredinger equation

$$\sum_{\alpha} \left(-\frac{\hbar^2}{2m_0} \Delta + V(r) \right) \Psi = E\Psi \quad (2.17)$$

comes to the view.

There are several ways to choose the appearance of the function $V(r)$ in solving a single-electron problem. The electron state is most applied, where all the ions of the lattice are in the field, their charge being compensated by the charge of the remaining electrons. Such a potential would be periodic, that is, it would have a three-dimensional periodicity of the grid.

In an ideal crystal, the atoms will be periodically arranged in space. Crystal

$$\mathbf{a}_n = n_1 \mathbf{a}_1 + n_2 \mathbf{a}_2 + n_3 \mathbf{a}_3 \quad (2.18)$$

When shifted to a vector, the atoms move on their own, where $\mathbf{a}_i$ - lattice constants, n_i - integer numbers. This means that the coordinate points r and $r + \mathbf{a}_n$ are equivalent from a physical point of view

$$V(\mathbf{r}) = V(\mathbf{r} + \mathbf{a}_n) \quad (2.19)$$

This expression implies that the potential area of the crystal is periodic.

If the specific value of energy in such a field is not decayed, then $\psi(r + \mathbf{a}_n)$ is also undecomposed and differs from $\psi(r)$ only to the constant

$$\psi(r + \mathbf{a}_n) = C\psi(r) \quad (2.20)$$

From Normalization Condition

$$|C|^2=1 \quad (2.21)$$

For being a permanent

$$C=\exp(ika_n) \quad (2.22)$$

will be equal to . Without it

$$|C|^2=|\exp(ika_n)|^2=|\cos ka_n + i \sin ka_n|^2 = \cos^2 ka_n + \sin^2 ka_n = 1$$

Here k - is a constant vector describing the quantum state of an electron in a crystal, called a wave vector. The number of wavelengths located at intervals between the modulus of the wave vector 2π is $k=2\pi/\lambda$ and is called the wave number.

(2.20) va (2.22) dan

$$\psi(r+a_n) = \exp(ika_n) \psi(r) \quad (2.23)$$

From here

$$\psi(r) = \exp(-ika_n) \psi(r+a_n) = \exp(ikr) U_k(r) \quad (2.24)$$

here

$$U_k(r) = \exp[-ik(r+a_n)] \psi(r+a_n) \quad (2.25)$$

The function $U_k(r)$ has the three-dimensional periodicity of the lattice

$$\begin{aligned} U_k(r+a_l) &= \exp[-ik(r+a_n+a_l)] \psi(r+a_n+a_l) = \\ &= \exp[-ik(r+a_n+a_l)] \exp(ika_l) \psi(r+a_n) = \exp[-ik(r+a_n)] \psi(r+a_n) = U_k(r) \end{aligned} \quad (2.26)$$

Thus, the stationary wave function of the electron in the periodic field of the crystal will depend on the wave vector

$$\Psi_k(\mathbf{r}) = e^{ikr} U_k(\mathbf{r}) \quad (2.27)$$

Where e^{ikr} - $\mathbf{k}$ is a flat wave running in the direction of the vector, $U_k(\mathbf{r})$ is a coordinate function that depends on the wave vector and has lattice periodicity. $\Psi_k(\mathbf{r})$ – Blox functions are called Blox functions. It follows that the wave function of an electron in the periodic field of the crystal is a function whose amplitude is modulated, has lattice periodicity, and runs in the direction of the wave vector.

If we express the Blox function (2.27) in (2.15)

$$H \Psi_k(\mathbf{r}) = E \Psi_k(\mathbf{r}) \quad (2.28)$$

This means that the energy of an electron will depend on the wave vector k in the crystal . This shows the law of dispersion for electrons in a crystal $E=E(k)$.

If the results obtained by the introduction of a self-adapted field depended on the finite value of the periodic potential, the solid-body theory would not exist until now, since neither its theory of the potential nor the results of experiment are known. But in order to get the fundamental results of the theory, it is enough not to have a finite value of a potential, but to know that it is periodicity in space and that this period is equal to the lattice constant.

2.4. Strongly bonded electron convergence

In describing the motion of electrons in a crystal lattice, the results of quantum mechanics provide results that are easily comparable to the quantum mechanical results for isolated atoms. As we know, the radiation spectrum of a free atom consists of discrete lines. Quantum theory shows that this means that in an isolated atom, electrons have discrete energy values.

One rule of physics is that a steady equilibrium state, that is, a state in which a system can be for a long time, is determined by a minimum of potential energy. In this context, all the electrons in the atom must be in the lowest energy state, i.e., at the energy level at the nearest state. But the laws of quantum mechanics put constraints on this state, which is shaped by Pauli's principle. According to Pauli's principle, only two electrons can exist at the same energy level, and they must have opposite spins.

Now, let us consider the interaction of many crystal-forming atoms. The energy levels of the internal electrons, that is, the electrons located close to the nuclei, are almost unchanged. This can be seen through X-ray private spectra, the appearance of which is virtually indistinguishable from any compound or substance by the state of the aggregate. However, the optical spectrum, that is, the transformation of electrons, caused by external electrons, changes dramatically.

If we assume that the kinetic energy of its electrons is much higher than the spatial change of its potential energy, then the periodic potential $V(r)$ is regarded as a small addition to the motion of free electrons. This approach generates an approximation called "near-free electrons," and this yields good results for some metals.

In the analysis of the physical properties of semiconductors, the strongly bonded electron approach becomes more accurate, in which the position of electrons in a crystal is perceived to be almost identical to their state in an isolated atom. But such an approach applies only to electrons in the deep energy layers of atoms, meaning they interact less with other nodes in the crystal lattice. Therefore, strong and weakly bonded electron approaches are not suitable for a complete calculation of the state of the valent electrons in the crystal. In other words, these approaches do not provide accurate results for calculating the energy spectrum of matter electrons, but they are well suited to explain the behavior of electrons in a crystal in a general way.

Suppose that $\mathbf{r}$ - is the radius vector of the electron, $\mathbf{R}_g$ - is the radius vector of the g -node of the crystal lattice. Then $\mathbf{r} - \mathbf{R}_g$ is the distance between the nucleus and the electron at the g -node, the wave function of the electron in the g -th atom is written as:

$$\Psi_g = \Psi_g(\mathbf{r} - \mathbf{R}_g) \quad (2.29)$$

In the strongly bonded electron approach, the position of an electron is perceived to be very similar to that of a crystal atom. Therefore, the wave function of an electron in a crystal can be expressed as a linear combination of the wave functions of an atom:

$$\Psi(\mathbf{r}) = \sum a_g \Psi_g \quad (2.30)$$

If we want to write this expression in the form of the Bloch wave function (2.27), then we need to put the following condition:

$$a_g = \exp(i\mathbf{k}\mathbf{R}_g) \quad (2.31)$$

With this, the wave function of an electron is written as follows:

$$\Psi(\mathbf{r}) = \sum \exp(i\mathbf{k}\mathbf{R}_g) \Psi_g(\mathbf{r} - \mathbf{R}_g) \quad (2.32)$$

It is considered correct to choose the wave function of the electron in the crystal in the form of (2.32). Because the wave function of an atom decreases exponentially, so the wave function of an electron located near the g-node is as follows:

$$\Psi(\mathbf{r}) \approx \text{const} \cdot \Psi_g(\mathbf{r} - \mathbf{R}_g) \quad (2.33)$$

That is, the wave function of an electron will be similar to the wave function of an atom at the g-node.

Now we put the expression (2.30) in Schrödinger's equation (2.15) and conclude as follows:

$$\sum a_g (H\Psi_g - E\Psi_g) = 0 \quad (2.34)$$

If $V(\mathbf{r})$ - is the potential energy of an electron in a crystal, then the Hamiltonian operator of the crystal is as follows:

$$\hat{H} = -\frac{\hbar^2}{2m_0} \Delta + V(\mathbf{r}) \quad (2.35)$$

Now, to determine the energy spectrum of the electrons in the crystal, let's assume that the energy spectrum of the insulated atom is known. The Hamiltonian of an isolated atom is expressed as:

$$\hat{H}_a = -\frac{\hbar^2}{2m_0} \Delta + U_a(\mathbf{r}) \quad (2.36)$$

where $U_a(\mathbf{r})$ - is the potential energy of the electron of the insulated atom. The Schrödinger equation is written for an isolated atom as follows:

$$H_a \Psi_a = E_a \Psi_a \quad (2.37)$$

To calculate the potential energy of an electron in a crystal, divide (2.36) by (2.35) and form the following expression:

$$H = H_a + W(\mathbf{r}) \quad (2.38)$$

here

$$W(\mathbf{r}) = V(\mathbf{r}) - U_a(\mathbf{r}) \quad (2.39)$$

The energy change of an electron in a crystal relative to an isolated atom.

Now, with the expression (2.38), we write the Schrödinger equation (2.34) as follows:

$$\sum a_g \{H_a + W(\mathbf{r}) \Psi_g - E \Psi_g\} = 0 \quad (2.40)$$

If we consider equation (2.37), (2.40) changes as follows:

$$\sum a_g \{W(\mathbf{r}) \Psi_g + (E_a - E) \Psi_g\} = 0 \quad (2.41)$$

By means of this equation, it is possible to study the energy spectrum of an electron in a crystal and its behavior.

(2.41) multiply the equation by Ψ_g^* and τ integrate it over the whole volume

$$\sum_g a_g \left\{ \int_{\tau} \Psi_g W(r) \Psi_g^* d\tau + (E_a - E) \int_{\tau} \Psi_g \Psi_g^* d\tau \right\} = 0 \quad (2.42)$$

We enter a mark

$$\int_{\tau} \Psi_g W(r) \Psi_g^* d\tau = A(R_g - R_{g'}) \quad (2.43)$$

This is an exchange integral, which refers to the degree to which atoms adhere to each other of wave functions and excitation energies, and the interaction determines the energy discharge.

$$\int_{\tau} \Psi_g \Psi_g^* d\tau = S(R_g - R_{g'}) \quad (2.44)$$

The integral refers to the degree of interlocking (closing) of the wave functions of atoms.

We rewrite (2.31), (2.43), and (2.44) to take into account the expressions (2.42)

$$\sum_g A(R_g - R_{g'}) e^{ikR_g} + (E_a - E) \sum_g S(R_g - R_{g'}) e^{ikR_g} = 0 \quad (2.45)$$

We multiply $e^{-ikR_{g'}}$ this equation by

$$\sum_g A(R_g - R_{g'}) e^{ik(R_g - R_{g'})} + (E_a - E) \sum_g S(R_g - R_{g'}) e^{ik(R_g - R_{g'})} = 0 \quad (2.46)$$

We introduce $q = R_g - R_{g'}$ definition

$$\sum_g A(q) e^{ikq} + (E_a - E) \sum_g S(q) e^{ik(q)} = 0 \quad (2.48)$$

From here we determine the specific values of the energy of the electrons in the crystal

$$E = E_a + \frac{\sum_g A(q) e^{ikq}}{\sum_g S(q) e^{ik(q)}} \quad (2.49)$$

This expression indicates that the energy of an electron in a crystal is equal to the sum of its specific surface energy in the isolated atom and the periodic function of the wave vector (k). The energy of an electron in a crystal does not come from

the state, but from the zones determined by the minimum and maximum values of the second hand.

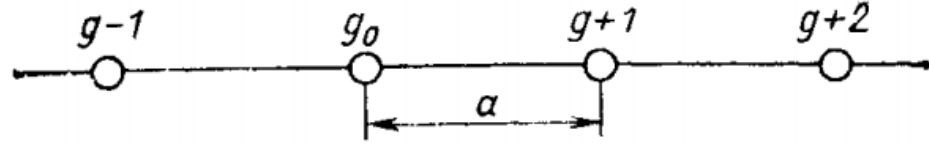


Fig. 2.2-One-dimensional lattice

We assume that the wave function of the atom in (2.28) is fast decaying, i.e., the interlocking of adjacent atoms is not accounted for. In this case, in (2.44), all the hads except $q=0$ become zero

$$S(q) = \begin{cases} 1, & q = 0 \\ 0, & q \neq 0 \end{cases} \quad (2.50)$$

From here

$$\sum_g S(q) e^{ikq} = 1 \quad (2.51)$$

It turns out that there is a lot of work to be done.

(2.49) when the image of the fraction in the expression is $q=0$

$$A(0) = \int_{\tau} \Psi_0 W(r) \Psi_0^* d\tau = \int_{\tau} |\Psi_0(r)|^2 W(r) d\tau = C \quad (2.52)$$

That is, $A(0)$ will not depend on k .

Even if the overlapping of the wave functions of neighboring atoms is not taken into account, the wave function near the zero node is partially compensated by the difference $W(r) = V(r) - U_a(r)$. For this reason, in the image of the expression fraction (2.49), it is also necessary to take into account the events adjacent to $q=0$.

If the wave function corresponds to the s-state, then the exchange integrals for all atoms will be the same. If the exchange integral is equal to A , then the energy of an electron in a crystal given (2.51) and (2.52)

$$E = E_0 + C + A \sum_g e^{ikq} \quad (2.53)$$

comes to the view.

In a simple cubic crystal lattice with a circuit, each atom is surrounded by six close atoms. If we direct the x, y, z axes to the edges of the cube, the adjacent nodes at the point $q=0$ will be located at the distances $q=+a$ and $q=-a$. So that:

$$\sum_{g_0} e^{ika} = e^{ik_x a} + e^{-ik_x a} + e^{ik_y a} + e^{-ik_y a} + e^{ik_z a} + e^{-ik_z a} = 2(\cos k_x a + \cos k_y a + \cos k_z a) \quad (2.54)$$

Thus, in a simple cubic lattice of a crystal, the energy of an electron is expressed as follows:

$$E = E_0 + C + 2A(\cos k_x a + \cos k_y a + \cos k_z a) \quad (2.55)$$

From this expression, several important conclusions can be drawn for the analysis of the electron's energy spectrum.

1. The E_a level of an isolated atom shifts to the amount C due to the interaction between atoms in the formation of a crystal lattice. The direction of displacement depends on the gesture of C .

2. The surface of the atom dissociates in the crystal lattice and forms a zone, inside which the energy of the electron changes periodically with respect to the components of the wave vector.

3. (2.55) The extreme values of the expression, when $\cos \cos(k_i a) = \pm 1$ ($i = x, y, z$), are as follows:

$$E_{max} = E_a + C + 6A,$$

$$E_{min} = E_a + C - 6A$$

In a simple cubic grid, the width of the energy zone is as follows:

$$E_{max} - E_{min} = 12A \quad (2.56)$$

that is, the exchange will depend on the value of the integral.

4. The energy level of each isolated atom is broken down into zones in the crystal. The magnitude of the interaction integral is determined by the interaction of electron clouds of neighboring atoms, that is, the stronger the wave functions of neighboring atoms interact, the larger they are, and consequently the width of the energy band is also large. Therefore, for higher atomic layers, a wider energy zone is formed due to the large interaction of electron wave functions (Figure 2.3).

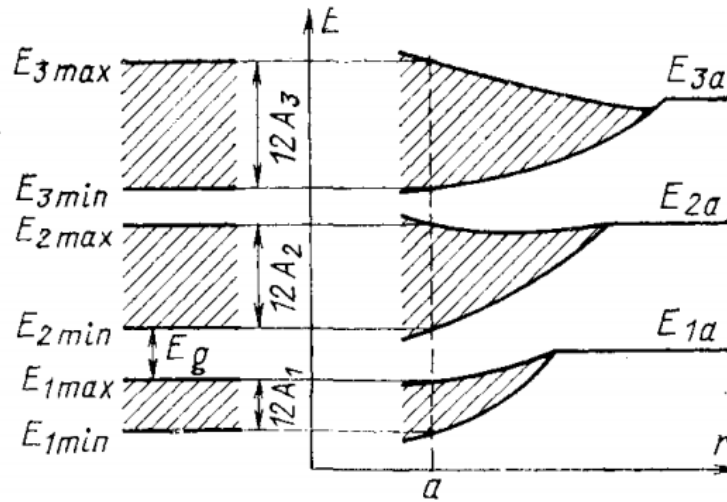


Fig. 2.3. Transformation of energy levels into zones when atoms are separated

5. Energy zones are generally separated by a prohibited energy range E_g , which are called prohibited zones (Fig. 2.3).

6. As energy increases, the width of the energy zones increases, while the width of the prohibited zones decreases (Fig. 2.3).

7. In an insulated atom, the level of E_a can be degraded. And in a crystal lattice, the rot can be partially or completely eliminated. Then the atomic level will decompose into several zones, the number of which will be equal to the decay rate. For example, for the p-state, the decay factor $g = 3$, since $g = 2l + 1$, where l - is the azimuth quantum number, and for the p-state $l = 1$. This means that three zones can be formed in the crystal from the p-state of the atom.

8. The energy of an electron in a crystal depends on the components k of the wave vector . It is a double function with respect to the wave vector, $E(k) = E(-k)$.

9. The temperature and pressures affecting the crystal cause the distance between the atoms to change, altering the closure region of the wave functions of the electrons, and thus changing the value of the substitution integral. This changes the width of the energy zones, and at the same time, the width of the forbidden zone between those zones also changes.

10. The strong bonding method does not apply to the outer valent electrons of the crystal because the width of the energy zone of the valent electrons will be equal to or greater than the distance between the energy levels in the insulated atom due to the large closure of the electron wave functions of neighboring atoms.

2.5. Number of electron states in the energy band

In order to calculate the number of states in energy, the conditions that must be introduced at the time of solving the Schrödinger equation are related to the constraints of the crystal, that is, it is necessary to take into account the conditions within the crystal boundaries. The boundary conditions of the crystal do not cause significant changes in the total energy of the electron states, since the number of nodes in true crystals is very large. Thus, the most convenient way to solve this problem is to introduce the cyclical conditions of Bor's - Karman. This approach considers the periodicity of the spatial properties of the electron in the crystal.

We consider the crystal in the form of a parallelepiped, with lengths L_x , L_y , and L_z on the x , y , and z axes. For a crystal having parameter a in a cubic lattice , there are the following relationships:

$$L_x = aN_x; L_y = aN_y; L_z = aN_z, \quad (2.57)$$

where N_x , N_y , N_z - is the number of atoms located in the crystal lattice on the x , y , z axes, respectively.

We require that the wave function Ψ have the same value on the opposite surfaces of the crystal. This condition imposes no physical constraints on the appearance of a wave function, only related to the boundaries of the crystal:

$$\Psi(x+L_x, y+L_y, z+L_z) = \Psi(x, y, z) \quad (2.58)$$

This condition basically represents the cyclic conditions of Bor's - Karman.

Considering the wave function of electrons in a crystal in the form of (2.27), we get the following equation:

$$\Psi(x+L_x, y+L_y, z+L_z) = U_k(r) e^{ikr} e^{ik_x L_x} e^{ik_y L_y} e^{ik_z L_z} = \Psi(x, y, z) \quad (2.59)$$

This condition must be fulfilled

$$e^{ik_x L_x} = e^{ik_y L_y} = e^{ik_z L_z} = 1 \quad (2.60)$$

Equality must be appropriate.

For this

$$k_x L_x = 2\pi n_1, \quad k_y L_y = 2\pi n_2, \quad k_z L_z = 2\pi n_3 \quad (2.61)$$

It should be, that is

$$k_x = \frac{2\pi}{L_x} n_1; \quad k_y = \frac{2\pi}{L_y} n_2; \quad k_z = \frac{2\pi}{L_z} n_3 \quad (2.62)$$

where n_1, n_2, n_3 - optional integers.

This means that the components of a wave vector are not continuous, but take on certain discrete values. Therefore, the energy of the electrons in the crystal is also quantized.

Considering the values of (2.32), considering a one-dimensional problem, we can write the wave function as follows:

$$\Psi_k = \sum_g e^{ik_x a g} \Psi_g = \sum_g e^{i \frac{2\pi n_x g}{N_x}} \Psi_g \quad (2.63)$$

Here it was assumed that $R=ag$, $L_x=aN_x$, the name g-atom.

From the expression (2.63) it follows that the wave function in $n_1=\pm N_x$ coincides with the wave function in $n_1=0$, and the wave function in $n_1=\pm(N_x+1)$ coincides with the wave function in $n_1=\pm 1$, and so on. This means that the component k_x will have a value N_x corresponding to different n_1 . In this case, since $E(k)=E(-k)$ n_1 accepts only the values 0, 1, 2, ..., N_x-1 , i.e.

$$0 < n_1 < N_x \quad \text{yoki} \quad -N_x/2 < n_1 < N_x/2 \quad (2.64)$$

From the expressions (2.57), (2.62), and (2.64), it follows that K lies at the following intervals

$$-\pi/a < k_x < \pi/a; \quad -\pi/a < k_y < \pi/a; \quad -\pi/a < k_z < \pi/a \quad (2.65)$$

Where k_x, k_y, k_z assume values N_x, N_y, N_z .

Consequently, in the permissible band of the crystal there are various energy states (energy levels) equal to the number of elementary cells in the crystal, a total of $N = N_x N_y N_z = L_x L_y L_z / a^3$ states.

According to quantum mechanics, the state of an electron in an atom is characterized by the prime quantum number n , the azimuthal quantum number l , the magnetic quantum number m , and the projection of spin from the z-axis s , that is,

four quantum numbers. Based on (2.62), the projections of the wave vector of the three quantum numbers should be k_x, k_y, k_z , and the fourth quantum number s_z . The projection of S can accept only two values: $+1/2$ and $-1/2$. Hence, in the case of (k_x, k_y, k_z) there can be no more than two electrons. Based on (2.55) (k_x, k_y, k_z) determines $E(k)$ for a given zone. Therefore, according to Pauli's principle, there can be no more than two electrons with oppositely directed spin electrons at each energy level defined by the wave vector. There are $2N$ quantum states at the unawakened atomic surface, respectively with N energy states and no more than $2N$ electrons, if the zone is awakened g times, it can contain $2gN$ electrons. It follows that the number of quantum states in the zone is equal to the total number of regions in the row of isolated atoms formed by this zone, that is, there is a conservation of the number of states in the formation of crystals from atoms.

Let's now calculate the average distance between adjacent energy levels in the permissible zone. Let us take the value of the lattice parameter $a = 4 \cdot 10^{-8}$ cm, the size of the unit cell is $a^3 = 64 \cdot 10^{-24}$ cm³. The number of elementary cells in a crystal with a unit volume of $V = 1$ cm³ is equal to the number of cases in the zone: $N = V/a^3 = 1/64 \cdot 10^{24} \sim 10^{22}$. In a zone with a width of 1 eV, the average distance between the layers is about 10^{22} eV, which means that the energy zone can be considered quasi-continuous.

2.6. Quasi-impulse

As we know, the state of a free-moving electron in space can be characterized by energy E and impulse P . In this, the relationship between energy and impulse is expressed as follows by the classical formula

$$E = P^2 / 2m_0 \quad (2.66)$$

De Brouil believed that a free electron with a mass of m_0 corresponds to a wave of length when v moves with velocity λ

$$\lambda = h/P = h/(m_0 v) \quad (2.67)$$

where h is Planck's constant.

To find

$$k = 2\pi/\lambda \quad (2.68)$$

is the impulse of a free electron, since it is

$$P = \hbar k \quad (2.69)$$

with the help of a formula, and the energy

$$E = \hbar^2 k^2 / 2m_0 \quad (2.70)$$

is determined using the formula.

For an electron moving in the periodic field of a crystal, the concept of a quasi-impulse $p = \hbar k$ can be introduced. Since k has discrete values, the quasi-

impulse also assumes discrete values. Quasi-impulse in a cubic lattice based on (2.65)

$$-\pi \hbar / a \leq p_i \leq \pi \hbar / a, \quad i=x, y, z \quad (2.71)$$

accepts values.

If we substitute p_i for k_i at (2.55), then for the electron energy in a cubic crystal

$$E = E_0 + C + 2A(\cos \frac{p_x a}{\hbar} + \cos \frac{p_y a}{\hbar} + \cos \frac{p_z a}{\hbar}) \quad (2.72)$$

We get the expression. From this expression it follows that the energy of an electron in a crystal is a double function of a quasi-impulse

$$E(\mathbf{p}) = E(-\mathbf{p}) \quad (2.73)$$

2.7. Brilluin zones

As shown in Section 2.5, in the energy zone of a crystal there are N energy states, the corresponding wave vector components are expressed as follows:

$$k_i = 2\pi n_j / L_i \quad (2.74)$$

and the components of quasi-impulse are:

$$p_i = 2\pi \hbar n_j / L_i \quad (2.75)$$

where $i=x, y, z; j = 1, 2, 3$ - integers.

For the cubic lattice, according to (2.65) and (2.71), the quasi-impulse components vary at the following limits:

$$-\pi/a \leq k_i \leq \pi/a, \quad -\pi \hbar / a \leq p_i \leq \pi \hbar / a \quad (2.76)$$

These values represent an area gathered around a point (000)(000) in the coordinate system of the quasipulse (p_x, p_y, p_z). This area is called the first or main Brillluien zone.

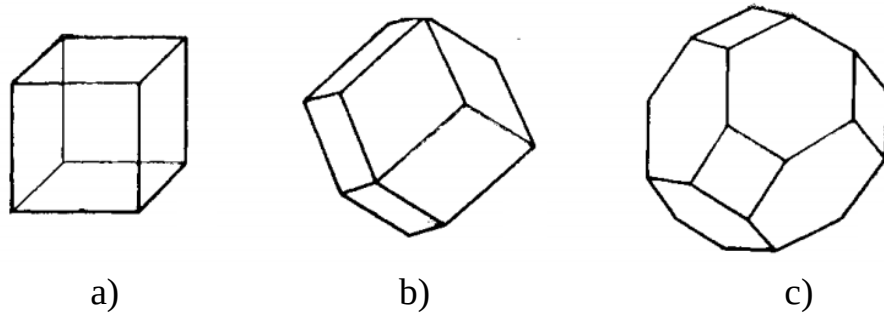


Fig. 2.4. The first Brillluien zone on cubic, volumetrically centered and surface-centered lattices

The first Brillouen zone for the cubic grid is cubic, the size of which is expressed as follows:

$$V_B^{(p)} = (2\pi \hbar / a)^3 = h^3 / a^3 \quad (2.77)$$

The Brillouen zone is represented as the first zone in the k-area, and this remains cubic for the grid as well

$$V_B^{(k)} = (2\pi / a)^3 = 8\pi^3 / a^3 \quad (2.78)$$

Fig. 2.5 depicts a section of the first Brillouen zone for a cubic lattice with a plane $p_z=0$. The corresponding points to the different states of the electrons are denoted by an "x". The center of the Brillouen zone (head of coordinates) denotes the state $n_1=n_2=n_3=p=0$. The distance between the nearest points in the Brillouen zone is equal to h/L based on (2.75). It follows from (2.63) that the dots marked with white circles in the figure denote the physique equivalent states located on opposite sides of the Brillouen zone. For example, points 1 and 2 represent equivalent states with the same wave function, in which the electrons in the holars will have the same energy.

Birinci Brillouen zonasini

$$\Delta V_B^{(p)} = (2\pi \hbar / L)^3 = h^3 / V \quad (2.79)$$

and

$$\Delta V_B^{(k)} = (2\pi / L)^3 = 8\pi^3 / V \quad (2.79a)$$

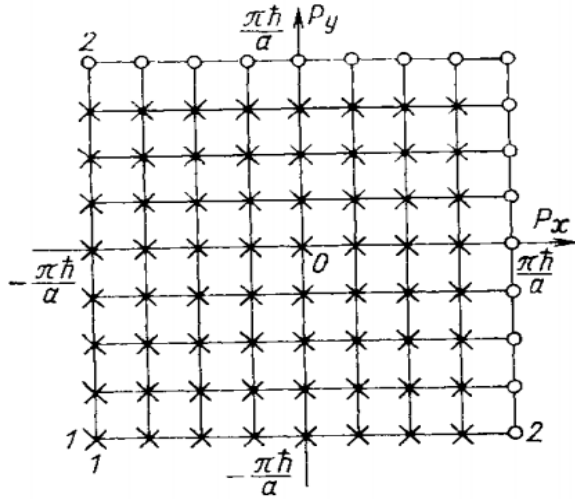


Fig. 2.5. Section of the first Brillouen zone for a simple cubic grid

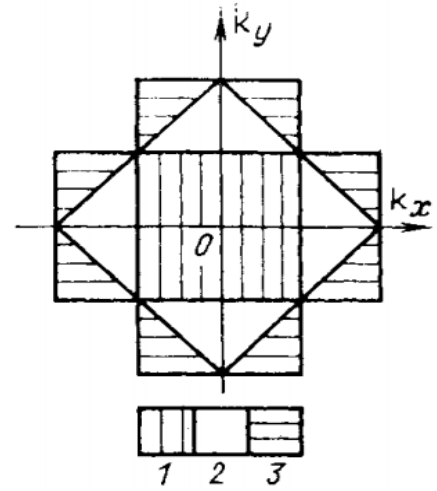


Fig. 2.6. The first three Brillouen zones for a regular cubic grid

can be divided into elementary cells, where $V=L^3=a^3N$ - is the crystal size, N - is the number of elementary cells in a crystal.

For a simple cubic crystal lattice, the size of the first Brillouen zone is $(h/a)^3$, and since the size of the elementary cell is h^3/a^3N , the number of elementary cells is equal to N , that is, the number of energetic states in the zone. But $2N$ electrons can

be settled in the energy zone. There may also be $2N$ electrons in the first Brillouin zone, and two electrons with opposite spin in each elementary cell.

The second and subsequent Brillouin zones, corresponding to the second and subsequent energy zones, will have more complex shapes, but their volumes will be constant. They are also composed of N elementary cells corresponding to the cells describing the equivalent states in the first zone. As an example, Fig. 2.6 depicts the first three Brillouin zones in k -space for a flat-curved lattice. Fig. 2.4-a for a simple cubic lattice, Fig. 2.4-b for a volumetric-centered lattice, and Fig. 2.4, c depicts the first Brillouin zones for a surface-centered lattice.

2.8. Filling the Valent Zone Electrons

Now let's look at a few examples of how the valent band is filled with electrons.

1. The zone is partially filled with electrons. If electrons are partially charged in the valent field, electrons can migrate to other empty energy levels under the influence of an external electric field. In this case, the presence of empty positions in the zone allows the electrons to move under the influence of the external electric field and to transport an electric charge. This means that in solids, for electric current to flow, there must be empty states in the valent zone.

Examples include monoatomic alkaline metals (*Li, Na, K, Rb, Cs*) and precious metals (*Cu, Ag, Au*) as electrical current good conductants. In alkaline metals, the valent electron is located in the "ns" state, while the remaining electrons are located in other states that are fully filled.

As an example, consider the element sodium, whose electrons are distributed as follows: $Na^{11} (1s^2 2s^2 2p^6 3s^1)$. During the formation of a crystal, the energy levels of sodium atoms decompose into zones. The $1s$, $2s$, and $2p$ layers are completely filled because they contain $2N$, $2N$, and $6N$ electrons. The electrons in the valent band are separated from the $3s$ state, of which N of the $2N$ states are empty. This means that even if the valent band in sodium is complete, under the influence of external energy, electrons pass into empty states and conduct electric current.

2. In the valent band, all possible electron states are occupied, but this band overlaps with the empty band occupied by electrons. In this case, the electrons of the valent zone pass to the energy layers of the free zone, making an electrical charge transfer. This situation is done in the elements of the second group. Alkaline earth elements have two valent electrons in the ns-valence state. For example, in $Mg^{12} (1s^2 2s^2 2p^6 3s^2)$, the $3s$ and $3p$ zones overlap, with the result that the electrons of the upper layers of the $3s$ -zone occupy the lower layers of the $3p$ zone.

3. The number of possible states in the stress band is equal to the number of valent electrons of the atoms forming the crystal. In this case, all states in the band are occupied by electrons, with two electrons with oppositely directed spin on each surface of the zone. Therefore, the external electric field cannot create a directed motion of such a set of electrons, since in a filled band, electrons can only exchange spaces. Therefore, such a crystal cannot conduct current, it is a dielectric;

As an example, consider the energy spectrum of a sodium chloride crystal. When $NaCl$ crystal is formed, the valent electron of sodium passes into the chlorine atom. The result is a negative chlorine ion and a positive sodium ion

$Na^+ (1s^2 2s^2 2p^6);$

$Cl^- (1s^2 2s^2 2p^6 3s^2 3p^6).$

The chlorine ion in the 3p-state has six electrons, which corresponds to the most stable energy state, that is, all the quantum states are occupied, and in each state there are two electrons with antiparallel spin. From this state, a zone is formed in the crystal where electrons occupy the entire 6N-state. It's going to be a valent zone. The sodium ion in the 3s state has no electron, and therefore a zone is formed in the crystal from this state where all 2N quantum states are free. It will be a permeability zone.

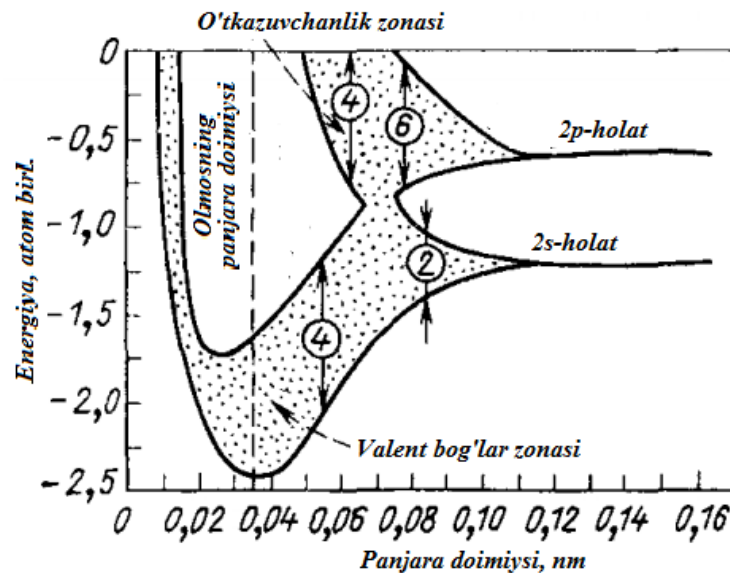


Fig. 2.7. The Transformation of Energy Stages in Diamond into Zones

Let us analyze the energy spectrum of crystals formed from the elements of group IV of the periodic table, which have a diamond-type crystal lattice. These include carbon (diamond), silicon, germanium and gray tin. The electronic structure of these atoms is such that (see Fig. 1.4 and 1.5) in the solid state, the four electrons of each atom are involved in the formation of covalent bonds. In this case, as shown in Fig. 2.7, the zones formed by ns - and np -states stick together to form a common zone where the number of cases is $8N$. As the interatomic distance decreases, this

zone splits into two zones, each having a quantum state of $4N$. In the lower band there are electron-filled states $4N$ - this is the stress band, and in the upper band the $4N$ quantum states are free - this is the conduction zone.

2.9. Dependence of electron energy on the wave vector at the bottom and ceiling of the zone

The energy of an electron in a simple cubic lattice, represented by Equation (2.55), is determined by the magnitude and gesture of the exchange integral. On the other hand, as can be seen from (2.43), the exchange integral A depends on the degree to which the atomic wave functions of the respective state adhere to each other, and the excitation energy depends on the gesture $W(r)$. Therefore, for the *s-state* with a spherical electron shell, $A < 0$. If we take into account the signal of the exchange integral, then the energy of the electron in the zone formed by the *s-state*

$$E = E' + 2|A|(\cos k_x a + \cos k_y a + \cos k_z a) \quad (2.80)$$

where $E' = Ea + C$ - is the energy of the *s-state*.

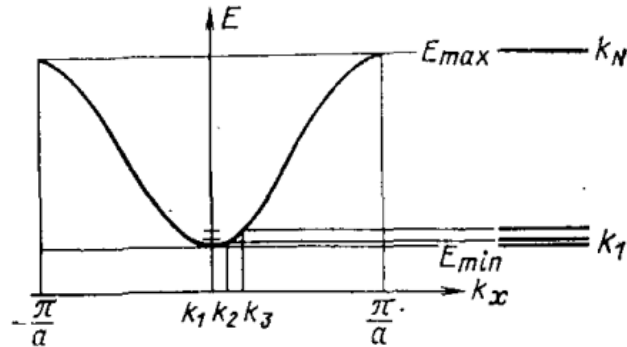


Fig. 2.8. Bonding $E(k_x)$ for the *s-state*

This energy link is shown in Fig. 2.8 of the variable wave vector change at the boundary of the First Brillouin Zone. The curve $E(k)$ is inverted for the k_y and k_z components of the *s-zone* and the wave vector.

In Fig. 2.8, the point energy $k_1=0$ corresponds to the minimum. The permissible levels of energy correspond to values $k=2\pi m/L_x$ $k_N=2\pi(N-1)/L_x$ - the point corresponds to the maximum-energy state.

$$xf(r)/r \quad (2.81)$$

For Ψ_p -states that are visible and partially cover the functions a of atoms, the atomic functions will have a basically opposite gesture.

It follows from Fig. 2.9 that the multiplication of the hidden parts of the same pointed functions is small and therefore can be considered $A > 0$. Then instead of (2.80)

$$E=E''+2|A|(\cos k_x a + \cos k_y a + \cos k_z a) \quad (2.82)$$

We obtain the expression.

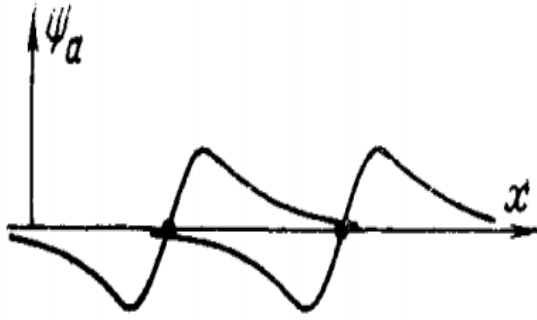


Fig. 2.9. Atomic functions for the p-state

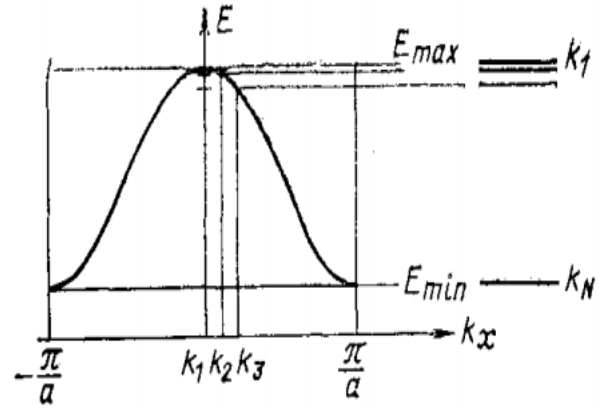


Fig. 2.10. Bonding $E(k_x)$ for state s.

Fig. 2.10 shows the dependence of energy on k_x for the zone formed from the p-state. Here $k_1=0$ corresponds to the maximum energy and k_N corresponds to the minimum energy.

For the zone formed from the d-state, the bond $E(k)$ is similar to the bond $E(k)$ for the zone formed from the s-state, and the graph is similar to the graph illustrated in Fig. 2.8.

Consider the expression (2.55) for the electron moving in the periodic field of a simple cubic lattice. If we define $E^a + C = 0$ as the initial enumeration system of energy, the electron energy in the zone generated by the s-state

$$E = -2A(\cos k_x a + \cos k_y a + \cos k_z a) \quad (2.83)$$

It is written in the form, where $A > 0$.

We enter a mark

$$E_{min} = -6A = E_c \quad (2.84)$$

$$E_{mak} = 6A = E_v \quad (2.85)$$

We analyze the expression around (2.83) point $k=0$. We spread the function $E(k)$ over the series and limit it to the first had because it is $ka \ll 1$

$$E = -2A \left(\left[1 - \frac{(k_x a)^2}{2} \right] + \left[1 - \frac{(k_y a)^2}{2} \right] + \left[1 - \frac{(k_z a)^2}{2} \right] \right) = -6A + Aa^2 k^2 \quad (2.86)$$

Thus, the electron energy at the bottom of the energy zone k will not depend on the direction of the wave vector k and will be proportional to k^2 .

(2.70) for the electron energy in the crystal, using (2.84) and (2.85) respectively

$$E = E_d + \frac{\hbar^2}{2m_n^*(d)} k^2 \quad (2.87)$$

We write expression, here

$$m_n^*(d) = \frac{\hbar^2}{2Aa^2} \quad (2.88)$$

m_n^* - The effective mass of the electron at the bottom of the energy zone.

Now we see the formula for cases near the upper limit of the zone (2.83). We introduce a definition $k' = \pi/a$. Since in this case there is a $kia \ll 1$ in the zone ceiling, let's spread out (2.85) into the row given (2.83),

$$\begin{aligned} E &= -2A[\cos\pi(-k) + \cos(\pi-k) + \cos(\pi-k)] = \\ &= 2A\left(\left[1 - \frac{(k'_x a)^2}{2}\right] + \left[1 - \frac{(k'_y a)^2}{2}\right] + \left[1 - \frac{(k'_z a)^2}{2}\right]\right) = 6A - Aa^2 k'^2 = E_m + \frac{\hbar^2}{2m_n^*(v)} k'^2 \quad (2.89) \end{aligned}$$

Thus, at the upper boundary of the energy band, the energy of the electron *will* be proportional to k'^2 and will not depend on k' , but the effective mass will be negative

$$m_n^*(v) = -\frac{\hbar^2}{2Aa^2} \quad (2.90)$$

Since the mass cannot be negative, the concept of a particle with a positive charge and a positive effective mass equal to the electron charge is introduced for the upper limit of the energy zone

$$m_p^*(v) = -m_n^*(v) = \frac{\hbar^2}{2Aa^2} \quad (2.91)$$

It follows that the effective mass of the pit at the upper boundary of the energy zone coincides with the effective mass of the electron at the bottom of this zone and is inversely proportional to the value of the exchange integral. Based on (2.56), the exchange integral determines the width of the energy zone. From this it follows that the effective mass of the charge carriers is inversely proportional to the width of the energy zone. This means that the higher the energy zone, the wider it is and the smaller the effective mass of charge carriers in it.

Fig. 2.11 shows the band $E(k) \pm \pi$ for the three adjacent zones of a one-dimensional lattice formed from the s-, p-, and d-states of atoms for the boundary change values of k . Each permissible zone is composed of N energy layers, with an average distance between levels of 10^{-22} eV. At points $k=0$ and $k=\pm\pi/a$, the electron enrichment assumes minimum and maximum values.

If we consider the space k_x, k_y, k_z or p_x, p_y, p_z ,

$$E(k) = \text{const} \text{ or } E(p) = \text{const} \quad (2.92)$$

The equation defines the constant energy surface or isoenergetic surface. According to (2.86) and (2.89), the energy of the electron and the lower and upper parts of the energy zone do not depend on the direction of the vector k , the isoenergetic surfaces

near the center of the Brillouin zone at $k = 0$ and $k = \pm p/a$ are spheres. As you move farther away from these points, the spheres begin to deform.

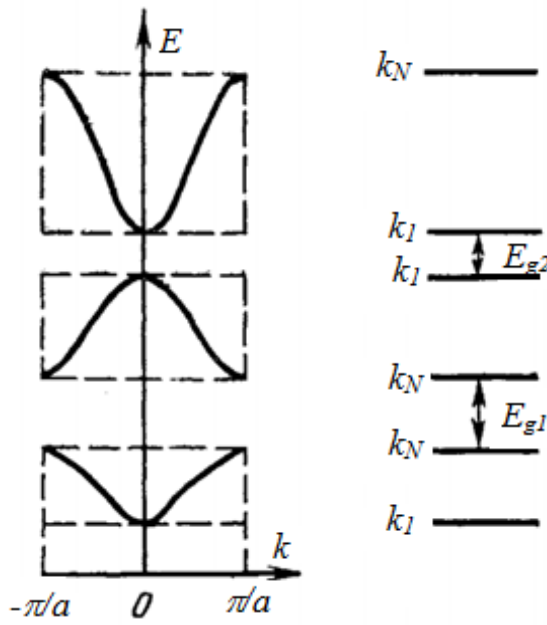


Fig. 2.11. $E(k)$ bonding for three adjacent zones formed from s-, p-, d- states of one-dimensional lattice atoms

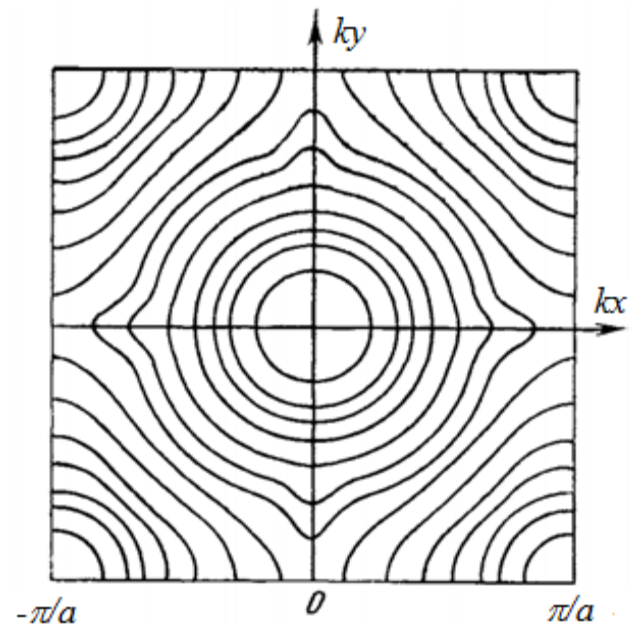


Fig. 2.12. Isoenergetic Lines for Square Flat Grid

Fig. 2.12 illustrates the appearance of isoenergetic surfaces (curves) for a square flat grid. Isoenergetic curves in the center and near the corners of the frame are circular in shape, and they break as they move away from these points.

2.10. Motion of electrons in a crystal under the influence of an external electric field

As determined, the energy spectrum of the electrons in the periodic field of the crystal has the form of quasi-continuous zones, separated by intervals of forbidden energy values (see Fig. 2.13).

We consider the motion of electrons in a crystal located in an external electric field directed along the z-axis $E = E_z = \text{const}$. In this case, the external force acts on the electron located in the periodic field of the F crystal, we consider the force F to be small enough and cannot cause the electron to pass from one energy zone to another, i.e.

$$Fa \ll E_g \quad (2.93)$$

Where a - is the constant of the E_g rid, the width of the elegant zone.

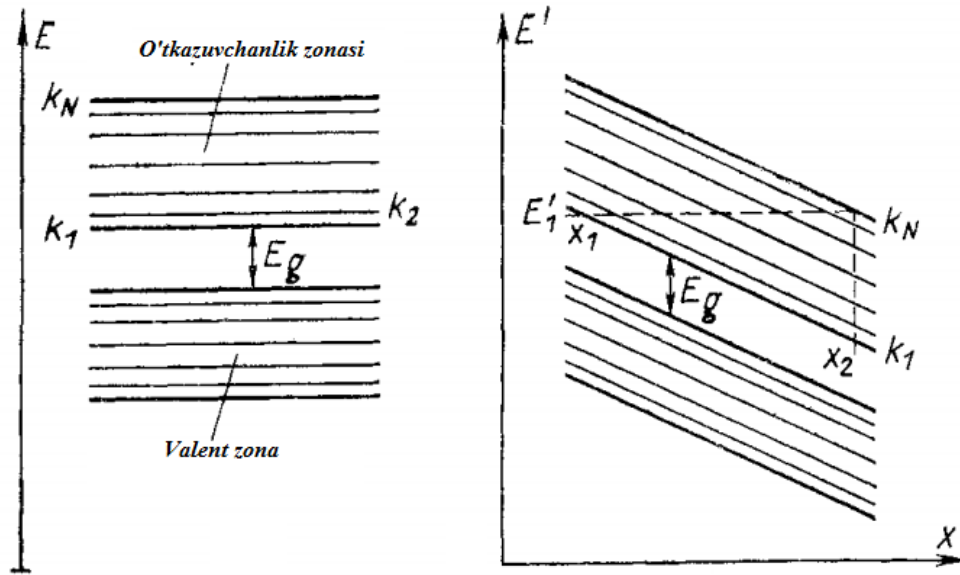


Fig. 2.13. Deflection of energy zones in the external electric field

Potential energy of an electron in an external electric field

$$U(r) = -eEx \quad (2.94)$$

is defined by expression.

The Schredinger equation for this case

$$\left(-\frac{\hbar^2}{2m_0} \Delta + V(r) + U(r) \right) \Psi = E' \Psi \quad (2.95)$$

will take on a look.

To simplify the problem of determining the specific values of energy, we assume that the external field in the field of the lattice constant is immutable $U = \text{const}$. In this case, you can use the Schrödinger equation

$$\left(-\frac{\hbar^2}{2m_0} \Delta + V(r) \right) \Psi = (E' - U) \Psi = E \Psi \quad (2.96)$$

Where E - is the electron energy in the absence of an external field. Energy of an electron in the external electric field

$$E' = E + U(r) \quad (2.97)$$

Without it, taking into account (2.94) (2.55) can be written as follows

$$E' = E_a + C + 2A(\cos k_x a + \cos k_y a + \cos k_z a) - eEx \quad (2.98)$$

From the expression (2.98) it follows that it is possible to obtain a graph of the coordinate dependence of the energy zones if all the values of the energy are given (Fig. 2.13). This graph denotes the deviation of the energy zones under the influence of the external electric field.

As you can see, the width of the zone for every value of x remains the same. Indeed, for the upper and lower tiers of the zone:

$$E_{\max} = E_a + C + 6A - eEx, \quad E_{\min} = E_a + C - 6A - eEx \quad (2.99)$$

$E_{max}-E_{min}=12A$ x does not depend on x and corresponds to the width of the zone for the stationary state.

In the absence of an external electric field, the electron will be at one of the levels of the energy zone and will have a certain value of the wave vector k , e.g., 2.13, as in Fig. 1. This means that in the stationary state, the electron has constant values of a quasi-impulse and a wave vector across the crystal at any distance.

When an electron acts in a crystal under the action of an external electric field, according to the law of conservation of energy, its total energy remains constant, i.e.

$$E' = E + U = const \quad (2.100)$$

However, the ratio between its kinetic and potential energies varies.

Fig. 2.13 – b shows the value of the total energy of an electron in a crystal in the presence of an external electric field by a puncture line parallel to the axis of the absiss. But in this case, as the electron moves across the crystal, it passes from one state of the energy band to another layer, as a result of which a change of the wave vector of the electron from k_1 to k_N occurs, and therefore its quasi-impulse also changes from p_1 to p_N . Jump to search Now an electron cannot move across the entire crystal, as shown in Fige 2.13 - b, but only moves in the range from x_1 to x_2 , that is,

$$\Delta x = x_2 - x_1 = (E_{max} - E_{min}) / eE \quad (2.101)$$

Thus, in the outer homogeneous electric field, the energy zones of the ideal crystal deviate, at which time the energy levels rise if $U > 0$ and decrease if $U < 0$. In this case, a redistribution of electrons occurs. Because they act in a given Δx field, the width of which is inversely proportional to the strength of the external electric field. In this field, when an electron is moving through the energy zone, it changes from state to layer and the wave vector and quasi-impulse of the electron change and so it becomes time-dependent, i.e.

$$k = k(t); P = P(t) \quad (2.102)$$

We define the law of change of the quasi-impulse and wave vector over time, that is, the law that describes the motion of an electron in a crystal in the presence of an external electric field.

Quantum mechanics already states that the motion of a free electron with a wave vector k can be described using Δx a wave packet consisting of a continuous variable plane wave superposition in the range of

$$k = k(t); P = P(t) \quad (2.102)$$

Wave function of a wave packet for a one-dimensional case

$$\Psi(x, t) = \int_{k-\Delta k}^{k+\Delta k} a(k) e^{i(kx - \omega t)} dk \quad (2.103)$$

It seems, where ω is the cyclic frequency.

The motion of the wave packet is characterized by the group velocity $v_{gr}=d\omega/dk$, which is equal to the velocity of some point in the packet, such as the maximum point. The coordinate of this maximum is determined from the condition $\omega x = dt/dk$

$$V=x/t=v_{gr} \quad (2.104)$$

That is, the average velocity of a free electron is equal to the group velocity of the wave packet.

The $E = \hbar\omega$ average velocity of a free circuit if a relation is used

$$v = d\omega / dk = \frac{dE}{\hbar dk} = \frac{dE}{dP} \quad (2.106)$$

is determined by the expression, where $P = \hbar k$ - free electron impulse.

The motion of an electron in a crystal is characterized by a wave function defined by a set of wave functions of an atom with different values of k . Since $k=2\pi n/aN$ ($n=0, 1, \dots, (N-1)$, $a=10^{-8}$ cm, $N=1022$ cm⁻³, $aN=1014$), the wave function can be considered as a set of flat waves with continuous variables of the wave vector. Therefore, the motion of an electron in a crystal can be described as a wave packet composed of Bloch functions. The expression (2.106) is also appropriate for the average velocity of the electron's motion in the crystal. Without three-dimensional

$$v = d\omega / dk = \frac{dE}{\hbar dk} = \frac{dE}{dP} \quad (2.108)$$

Here $P = \hbar k$ - quasi-impulse.

Thus, the average velocity of motion of an electron in a crystal is equal to the derivative of the energy in terms of quasiimpulse.

In the strong bond approximation, we calculate the average quantummechanical velocity of an electron in a simple cubic lattice.

From equations (2.83) and (2.108)

$$v = \frac{2aA}{\hbar} (i \sin k_x a + j \sin k_y a + k \sin k_z a) = \frac{2aA}{\hbar} (i \sin \frac{p_x a}{\hbar} + j \sin \frac{p_y a}{\hbar} + k \sin \frac{p_z a}{\hbar}) \quad (2.109)$$

We take the parity, where i, j, k - Cartesian coordinate backs.

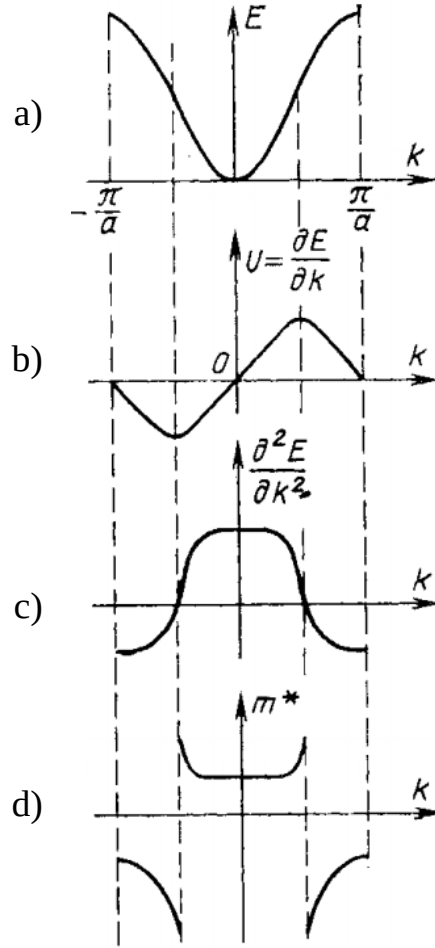


Fig. 2.14. Dependence of energy, velocity, $\partial^2 E / \partial k^2$, and effective mass m^* on k for the cubic lattice

It follows from the expression (2.109) that the average speed of the electron's motion in the crystal depends on the modulus and direction of the wave vector (quasi-impulse).

In the one-dimensional case, $k_y = k_z = 0$, $k_x = k$ and

$$E = -2A \cos ka; \quad (2.110)$$

$$v = \frac{2aA}{\hbar} \sin ka \quad (2.111)$$

Fig. 2.14-a and b illustrate the graphs of these links. From the figures it follows that at the boundaries of the zones ($k = \pm\pi/a$) $E(k)$ reaches extreme values. This means that the electron velocity at the extreme points is zero. And inside the zone reaches its maximum value. Mean quantummechanical velocity of the electron across the zone

$$v_{o'rt} = \int_{-\pi/a}^{\pi/a} v(k) dk = \frac{2aA}{\hbar} \int_{-\pi/a}^{\pi/a} \sin kadk = 0 \quad (2.112)$$

This means that in the absence of an electric field, the mean quantummechanical velocity of an electron across the crystal zone is zero.

Suppose an external force is acting on an electron in a crystal. Based on the law of conservation of energy for a one-dimensional case

$$\frac{dE}{dt} = Fv \quad (2.113)$$

that is, the rate of change of energy in a unit of time is equal to the work performed by external forces

$$\frac{dE}{dt} = \frac{dE}{dp} \frac{dp}{dt} \quad (2.114)$$

(2.113), (2.114) and (2.107)

$$\frac{dE}{dt} = F \quad (2.115)$$

origin.

This equation is identical in shape to Newton's equation of motion: the rate of change of a quasipulse is proportional to the force exerted. Consequently, in the periodic field of the $V(r)$ crystal, the electron quasi-impulse p in the equation of motion (2.115) performs the role of the free electron impulse.

Let the external force F be dependent on time. In this case, the motion of the electron in p -space

$$p(t) = p_0 + Ft \quad (2.116)$$

It is characterized by an equation and the trajectory of motion is a straight line, the direction of which is determined by the direction of an external force. But *in p -space* (or k -space), an electron moving from one isoenergetic surface to another (see Fig. 2.12). Therefore, the external force F changes not only the quasi-impulse but also the energy of the electron.

Now we look at how the P impulse of the electron in the crystal changes in the absence of an external electric field. In a crystal with an ideal structure, a rigid periodic field, an electron acts in a uniform level of energy field. Since the electron impulse is constant, $dp/dt = 0$. But by the lattice field, the crystalline electron is affected by *the force F* and determines the change in the impulse, i.e.

$$dP/dt = F \quad (2.117)$$

Since F is a periodic function equal to the lattice circuit, the impulse of an electron moving in a crystal is not constant, but changes periodically.

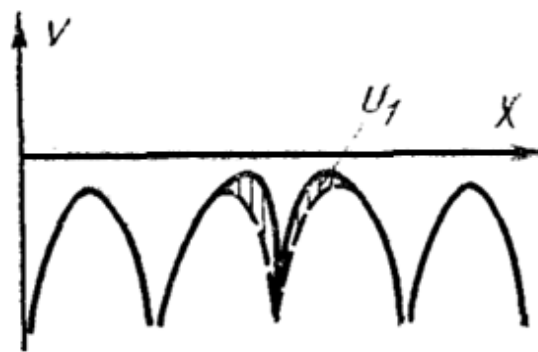
When there is an external electric field, the electron is affected by both the external force F and the force F_{kr} caused by the action of the lattice. Therefore, the change of the electron impulse over time is determined by both the influencing and external force from the crystal area, i.e.

$$dP/dt = F_{kr} + F \quad (2.118)$$

Thus, if the structure of the crystal lattice is ideal, then its potential area is $V(r)$ solid-periodic. In the periodic lattice field, the electron moves through the entire

crystal with a constant quasi-impulse and therefore constant energy. In this case, the electron moves along the crystal and remains at the same level of the energy zone. But since $p=\text{const}$ and therefore $k=\text{const}$, the electron velocity according to (2.111) is constant $v=\text{const}$, so its acceleration is zero. This means that in the circuit field of the grid, the electron moves without acceleration. In other words, in a solid circuit field, an electron moves like a free particle without resistance, without scattering.

If a crystal with ideal structure is placed in an external electric field, then according to (2.115), the motion of an electron under the influence of an external force F will be similar to the motion of a free particle. Only in this case, an electron moving in a limited region of the crystal (2.101) in the energy zone shifts from state to state and its energy changes.



2.15-Fig. 2.15. Distorted Potential of a One-Dimensional Grid

Now let's consider the case in which the crystal lattice area is not rigid. Let the grid area be similar to the one shown in Fig. 2.15. For the part of the structure where the ideality of the crystal structure is broken, the lattice area can be shown as a periodic $V(r)$ area, on top of which some additional force field U_1 with lattice periodicity is placed. But according to (2.115), the quasi-impulse of the electron must change when U_1 is exposed to a non-periodic part of the potential field. It follows that any violation of the ideal periodicity of the lattice potential field leads to a change in the electron quasi-impulse. This means that any violation of the ideal of the lattice structure will cause electron wave propagation.

In real crystals, the potential field is not solid. It is disrupted by thermal vibrations of the lattice atoms and the presence of various types of defects in it. The propagation of electron waves by the inhomogeneity of the lattice structure is a physical cause that determines its resistance when an electric current passes through a substance.

2.11. Effective mass of charge carriers

Suppose that a free electron with mass m is located in a homogeneous electric field whose voltage is E . An electron is exposed to a force F , under the influence of which the electron receives acceleration

$$a = -eE/m_0 \quad (2.119)$$

Given (2.107) and (2.115) for the electron in the crystal located in the outer field

$$a = \frac{\partial^2 E}{\partial p^2} F \quad (2.120)$$

It can be written. For a three-dimensional case

$$\begin{cases} a_x = \frac{\partial^2 E}{\partial p_x^2} F_x + \frac{\partial^2 E}{\partial p_x \partial p_y} F_y + \frac{\partial^2 E}{\partial p_x \partial p_z} F_z \\ a_y = \frac{\partial^2 E}{\partial p_y \partial p_x} F_x + \frac{\partial^2 E}{\partial p_y^2} F_y + \frac{\partial^2 E}{\partial p_y \partial p_z} F_z \\ a_z = \frac{\partial^2 E}{\partial p_z \partial p_x} F_x + \frac{\partial^2 E}{\partial p_z \partial p_y} F_y + \frac{\partial^2 E}{\partial p_z^2} F_z \end{cases} \quad (2.121)$$

In this case, the direction of the acceleration vector does not coincide with the direction of the external force vector.

Coupling an acceleration vector and an external force vector

$$\frac{\partial^2 E}{\partial p_i \partial p_j} = \frac{1}{m_{ij}} \quad (2.122)$$

Magnitudes are called elements of the second color tensor or the inverse effective mass tensor.

We spread the expression of electron energy in a crystal around the extremal point k_0 into a series of degrees of $k-k_0$ and are satisfied with second-order hads

$$E(k) = E(k_0) + \sum_{ij} \left(\frac{\partial^2 E}{\partial p_i \partial p_j} \right)_{k_0} (k_i - k_{i0})(k_j - k_{j0}) \quad (2.123)$$

Consider $i, j=x, y, z$. (2.122)

$$E(k) = E(k_0) + \frac{\hbar^2}{2} \sum_{ij} \frac{1}{m_{ij}} (k_i - k_{i0})(k_j - k_{j0}) \quad (2.125)$$

or

$$E(p) = E(p_0) + \frac{1}{2} \sum_{ij} \frac{1}{m_{ij}} (p_i - p_{i0})(p_j - p_{j0}) \quad (2.126)$$

The velocity of motion of an electron in a crystal around a point k_0

$$v_i = \frac{\partial E}{\partial p_i} = \sum_{ij} \frac{p_i - p_{i0}}{m_{ij}} \quad (2.127)$$

would be equal to, for example,

$$v_x = \frac{p_x - p_{x0}}{m_{xx}} + \frac{p_y - p_{y0}}{m_{xy}} + \frac{p_z - p_{z0}}{m_{xz}}$$

or in vector form

$$\mathbf{v} = \frac{\mathbf{p} - \mathbf{p}_0}{m} \quad (2.128)$$

This means that the velocity of motion of an electron in a crystal is equal to the scalar product of the inverse effective mass tensor with a quasi-impulse.

(2.122) Since the tensor is symmetric with respect to the prime diagonal, it is possible to bring the elements of the prime diagonal to a symmetric tensor when they are non-zero. In this case, the inverse effective mass tensor is called the inverse tensor of the effective mass tensor.

Expression given effective mass tensor (2.125)

$$E(k) = E(k_0) + \frac{\hbar^2(k_x - k_{x0})^2}{2m_1} + \frac{\hbar^2(k_y - k_{y0})^2}{2m_2} + \frac{\hbar^2(k_z - k_{z0})^2}{2m_3} \quad (2.131)$$

or

$$E(p) = E(p_0) + \frac{(p_x - p_{x0})^2}{2m_1} + \frac{(p_y - p_{y0})^2}{2m_2} + \frac{(p_z - p_{z0})^2}{2m_3} \quad (2.132)$$

is written in the form. Velocity is

$$\mathbf{v}_i = \partial E / \partial \mathbf{p}_i \quad (2.133)$$

is defined by expression.

(2.132) It follows from the expression that surfaces of the same energy in p-space are ellipsoids around the extreme point of energy, and m_1, m_2, m_3 will be the head axes of the ellipsoids.

Ellipsoid tenglamasini

$$\frac{(p_x - p_{x0})^2}{a^2} + \frac{(p_y - p_{y0})^2}{b^2} + \frac{(p_z - p_{z0})^2}{c^2} = 1 \quad (2.134)$$

We write in the form, where a, b, c - are the half axes of the ellipsoid in the directions x, y, z

$$a^2 = 2m_1[E(p) - E(p_0)]; \quad b^2 = 2m_2[E(p) - E(p_0)]; \quad c^2 = 2m_3[E(p) - E(p_0)]. \quad (2.135)$$

Since the gradient is directed to the surface and the electron velocity is the energy gradient in quasi-impular space, it is perpendicular to the isoenergetic surface. Therefore, for ellipsoidal surfaces of energy, the directions of velocity and quasi-impulse may not be matched. Their orientation coincides along the ellipsoid axes. In this case, if $E(p_0) = E_0$

$$\mathbf{v} = \frac{\mathbf{p}_i - \mathbf{p}_{i0}}{m_i} = \sqrt{\frac{2(E - E_0)}{m_i}} \quad (2.137)$$

From this expression, the velocity of an electron along the ellipsoid axis at the same energy is inversely proportional to the corresponding component of the effective mass. But the effective mass based on (2.135) determines the axes of the ellipsoid. Therefore, the longer the ellipsoid axes, the smaller the electron velocity.

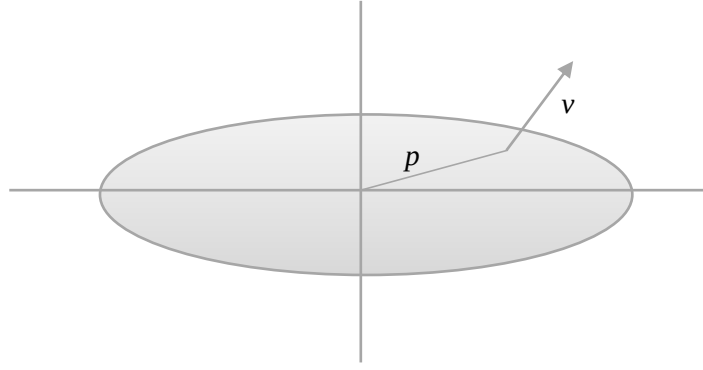


Fig. 2.16. Velocity and quasi-impulse of an electron on an ellipsoidal energy surface

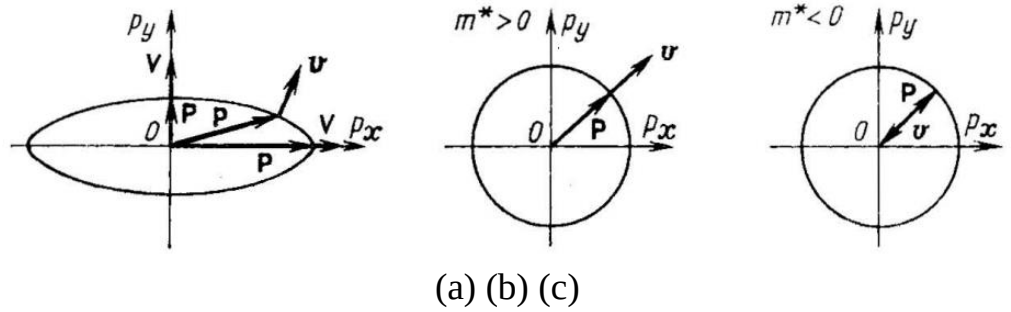


Fig. 2.17. Radius vector and direction of normal on ellipsoidal, spherical ($m^* > 0$), spherical ($m^* < 0$) surfaces of energy

If the two axes of an ellipsoid are equal, then the effective masses are also equal $a=b$, $m_1=m_2$. Isoenergetic surfaces then

$$E(p) = E(p_0) + \frac{(p_x - p_{x0})^2}{2m_1} + \frac{(p_y - p_{y0})^2}{2m_2} + \frac{(p_z - p_{z0})^2}{2m_3} \quad (2.138)$$

will appear to be.

The isoenergetic surface described by Equation (2.138) is the rotational ellipsoid corresponding to the axis of rotation p_z . Since m_2 is the effective mass along the axis of symmetry, i.e., along the axis of rotation of the ellipsoid, it is called the longitudinal effective mass of the axis and is denoted by m_1 . Since the symmetry m_1-m_2 is transverse to the axis, it is called the transverse effective mass. If $m_1 < m_2$, then the ellipsoid of rotation will be elongated along the axis of rotation, and the greater the ratio of effective masses m_2/m_1 , then the ellipsoid will be elongated. Otherwise it will be compacted.

For crystals with cubic symmetry, all major axes are equivalent. In this case, $m_1=m_2=m_3=m^*$ and the effective mass tensor becomes a scalar [zero-degree tensor]. In this case, isoenergetic surfaces

$$E(p) = E(p_0) + \frac{p^2}{2m^*} = \text{const} \quad (2.139)$$

It is characterized by an equation and has a spherical appearance (Figure 2.17- b). The effective mass expression looks like this:

$$\frac{1}{m^*} = \frac{\partial^2 E}{\partial p^2} = \frac{\partial^2 E}{\hbar^2 \partial k^2} \quad (2.140)$$

Therefore, for a cubic symmetry crystal, the effective mass of the charge carriers is inversely proportional to the curvature of the isoenergetic surfaces in k -space (or p -phase). When the electron is near the energy minimum, that is, near the bottom of the conduction zone (Fig. 2.14-a).

$$\frac{\partial^2 E}{\partial k^2} > 0, \quad m^* = \text{const} > 0 \quad (2.141)$$

That is, electrons behave like negatively charged particles with a positive effective mass (Fig 2.14.d).

In this case, in accordance with the phrases (2.120) and (2.140), we get:

$$F = ma \quad (2.142)$$

That is, the direction of acceleration is in the direction of the external force and is therefore opposite to the direction of the external electric field.

Since the acceleration is $a = dv/dt$, $F = dp/dt$, the quasi-impulse of an electron in a cubic crystal is equal to:

$$p = m^*v, \quad (2.143)$$

that is, the velocity of electron motion is in the quasi-impulse direction (Fig. 2.17-b).

Consequently, the motion of an electron at the bottom of the energy zone of the cubic crystal under the influence of an external electric field is similar to the motion of a free particle whose mass is equal to m^* .

From (2.142) it follows that the acceleration of the electron in the crystal is only influenced by an external force. The effect of the lattice electric field is manifested in the fact that the motion of an electron under the influence of an external force is determined not by its usual mass m_0 , but by its effective mass m^* .

Near the energy maximum, i.e., near the upper part of the tensile zone (Fig. 2.14.a),

$$\frac{\partial^2 E}{\partial k^2} < 0, \quad m^* = \text{const} < 0 \quad (2.144)$$

(Fig. 2.14, d), and therefore velocity is directed against the quasi-impulse (Fig. 2.17-c.). It follows from (2.144) that the direction of acceleration of an electron located at the upper edge of the zone is opposite to the direction of the external force acting on it:

$$a = -F/m^* \quad (2.145)$$

But in a homogenous electric field, the electron

$$F = -eE, \quad (2.146)$$

The force acts, so that for an electron located close to the ceiling of the energy zone where $m^* < 0$, the acceleration by the external electric field is as follows:

$$a = eE/m^* \quad (2.147)$$

That is, the acceleration of the charge carrier will be the same as the direction of the electric field.

Such a charge carrier located near the top of the stress zone behaves like a particle with a positive charge and a positive effective mass, and is called a pit. When moving away from the edges of the zone, the dependence $E(k)$ deviates from the quadratic law. At the turning points $k = \pm\pi/2a$ of the curve (Fig. 2.14)

$$\frac{\partial^2 E}{\partial k^2} = 0, \quad m^* \rightarrow \infty$$

It follows that in order to describe the motion of the charge carriers in the crystal under the influence of an external electric field (2.115), the relation can only be applied to the charge carriers located near the bottom or ceiling of the energy zone. In practice, the two states of energy zone replenishment are considered the most important.

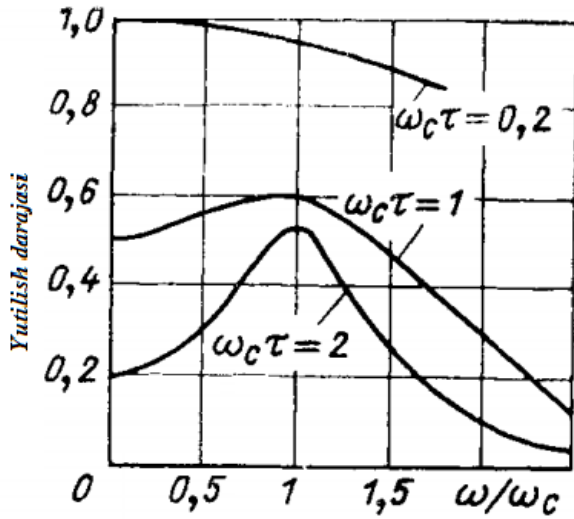


Fig. 2.18. Absorption in cyclotron resonance depending on constant magnetic field strength

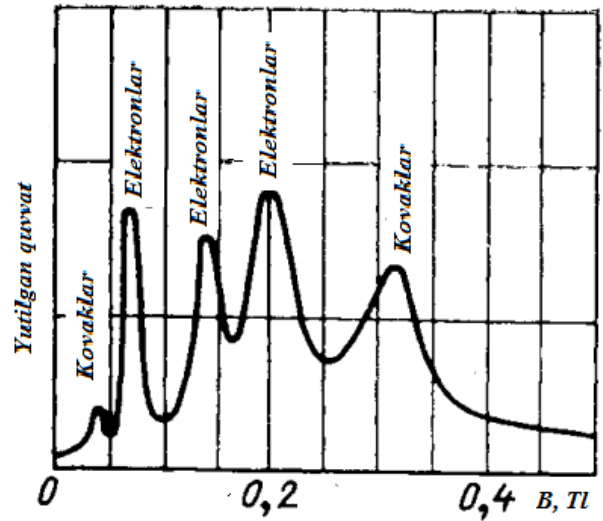


Fig. 2.19. Absorption in Germanic cyclotron resonance

1. The zone is almost empty. In this case, the number of quantum states in which the number of electrons in the band n is much less than $2N$. Electrons that tend to occupy lower energy levels are located at the bottom of the zone. In this case, $m_n^* = \text{const} > 0$.

2. The zone is almost filled with electrons. In such a zone there are a certain number of free quantum states. Electron-free states are located in the highest part of the zone, that is, the hives are located at the top of the zone, and their energy is

minimal. Therefore, the energy of the holes is measured in the opposite direction to the energy of the electron. In this case, $m_p^* = \text{const} > 0$.

At absolute zero temperature, the valence zone of the semiconductor is filled with electrons, and the conduction zone is free. With a temperature increase, a small section of electrons pass from the valence zone into the conduction zone due to thermal excitation. In this case, the condition considered above will be performed: the presence of a small number of electrons at the bottom of the conduction zone and holes at the upper edge of the valence zone. Since the width of the valence band is greater than the width of the conduction band, and since the value of the effective mass is inversely proportional to the width of the energy band, the effective mass of the hole must be greater than the effective mass of the electron, i.e.

$$m_p^* > m_n^*$$

This means that the electron and the hole differ not only in their charge, but also in their effective mass.

2.12. Cyclotron resonance

There are several experimental methods for determining the effective mass of charge carriers. One of them is the cyclotron resonance method. In this case, the semiconductor is placed in a constant magnetic field with $B = \text{const}$. A circuit moving at a velocity v is affected by the Lorentz force

$$F = -e[vB] \quad (2.148)$$

Under the action of this force, the electron moves in a circle in a plane perpendicular to the magnetic field. The cyclic frequency, called the cyclotron frequency, is determined by the equation of centrifugal force and Lorentz force

$$m^*v^2/r = eBv \quad (2.149)$$

$$\omega = v/r = eB/m^* \quad (2.150)$$

It follows that the frequency of the cyclotron does not depend on the velocity and radius of the circle. If now a weak high-frequency electric field oscillating in a plane perpendicular to B is formed in a semiconductor, then when approaching the field frequency, the electromagnetic field absorbs the energy vigorously. This resonance absorption phenomenon is commonly referred to as cyclotron resonance. The resonance frequency of absorption is determined by measurements, followed by (2.150) which is the effective mass of the electron (or hole). If this amount is anisotropic, it is possible to study the ellipsoid shape of the same energy by changing the direction of the magnetic and electric fields with respect to the crystallographic axes of the crystal.

If the free travel time of the electron is large enough, an experiment can be performed to determine the frequency of resonant absorption:

$$\tau \gg \omega$$

That τ is, over time, the circuit should be able to spin at least once.

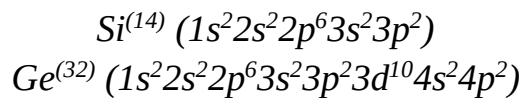
The more the electron rotates, the sharper and sharper the resonance. Fig. 2.18 $\omega_c \tau$ shows the absorption curves of the electromagnetic field for the three values of c (0,2; 1 and 2):

The greater the free flow time of the charge carrier, the fewer structural flaws there will be in the crystal. Therefore, cyclotron resonance experiments must be conducted at very low temperatures (typically 4.2 K, the temperature of liquid helium) and in highly pure crystals. Typically, in experiments, the frequency of radio radiation is set and resonance is achieved by changing the magnetic field induction. Fig. 2.19 shows the absorption spectrum of the cyclotron resonance in Germany.

2.13. The bandwidth structure of some semiconductors

Calculating the zonal structure of semiconductors is a complex matter. For germanium and silicon, the bonding of $E(k)$ would be the 146th degree equation. The energy is therefore calculated only for some symmetric points, which greatly simplifies the calculation. However, even in this case, it will be necessary to solve the 16-degree equation in a finite way. For intermediate points, the Brilluin zone is determined by interpolation. The results of theoretical calculations are compared with experimental data, including in terms of cyclotron resonance.

The electronic structure in silicon and german is as follows:



Their last shell is unfilled, since there are two electrons with parallel spin in the p-state, the conductivity zone and the stress zone of silicon and germanium contain the p-state (see Fig. 2.1), for which the excitation is lost in the crystal, each representing the superposition of three different zones. In Fig. 2.21, they show three $B(k)$ represented by branches. This connection is not the same for different crystallographic directions.

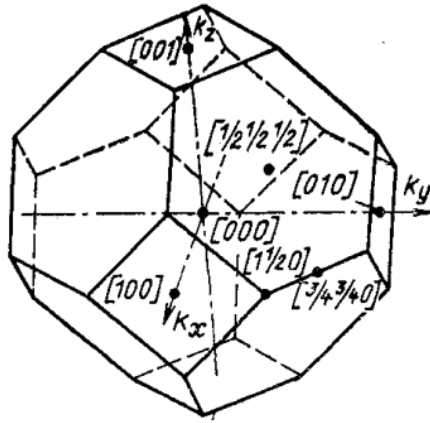


Fig. 2.20. The First Brillouin Zone for Diamond-Type Semiconductor

In silicon and germanium, one of the $E(k)$ horns of the permeability zone (Fig. 2.21) lies significantly lower than the others. The state of the absolute minimum of energy determines the bottom of the permeability zone. Energy Ministries are also called valleys.

In Germany, the absolute minimum of the conductivity zone $[111]$ is located along the direction of the axes (Fig. 2.21), so there are eight equivalent energy minimums, the coordinates of which are visible in *units* $2\pi/a$ ($1/2, 1/2, 1/2$), which lie at the boundary of the Brillouin zone. That is, the energetic surfaces near absolute energy minimums (or valleys) are ellipsoids relative to the large semi-axis. In German, it $[111]$ corresponds to the direction of the axes, and the first Brillouin zone is half of each energy ellipsoid, so in Germany there are only four fully energetic ellipsoids instead of eight (Fig. 2.22).

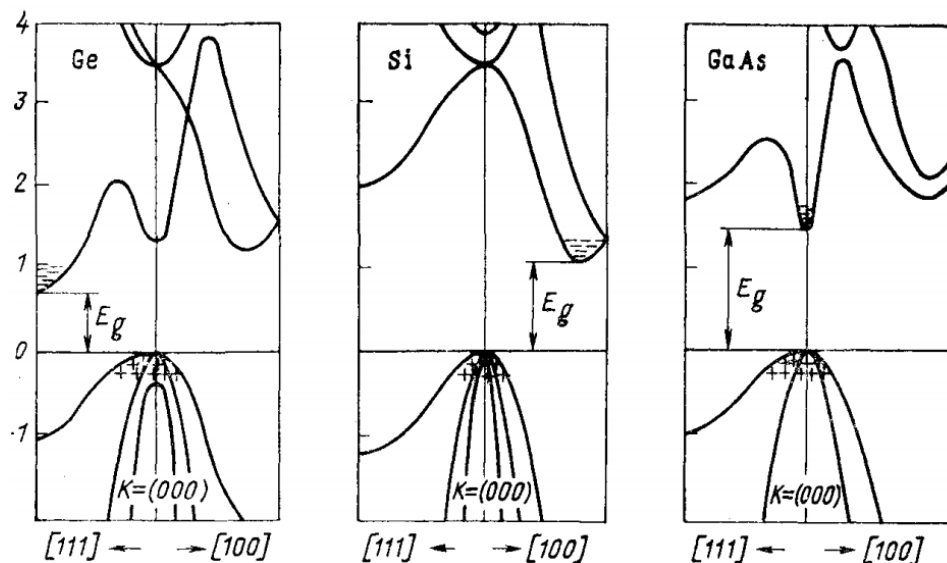


Fig. 2.21. Energy zones of germanium, silicon and gallium arsenide

The absolute minimum of the conductivity zone in silicon lies in the direction of the [100] axes, not far from the boundary of the Brillouin zone (Figure 2.21). This means that silicon has a minimum of six equivalent energies, and this means that in the first Brillouin zone of silicon [100] there are six ellipsoidal surfaces with constant energies stretched along the axes (Figure 2.22). The centers of ellipsoids are about three-quarters of a mile from the center of the Brillouin zone.

In gallium arsenide (see Figure 2.21) and indium antimonide (Fig. 2.25), the absolute minimum of the conductivity zone *lies at* $k = 0$, so that their constant-energy surfaces are spheres whose center is located in the center of the Brillouin zone, and the effective masses of electrons are $0.068m_0$ and $0.010m_0$, respectively.

Dependence of electron energy on a wave vector k near the absolute minimum of the conductivity zone in germanium and silicon

$$E(k) = E(k_0) + \frac{\hbar^2[(k_x - k_{x0})^2 + (k_y - k_{y0})^2]}{2m_l} + \frac{\hbar^2(k_z - k_{z0})^2}{2m_s} \quad (2.152)$$

It is represented by the formula.

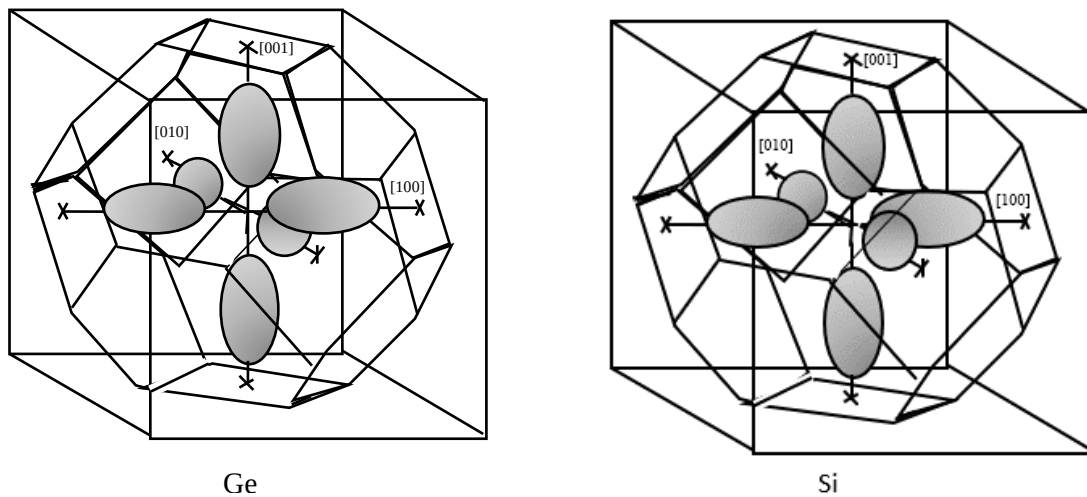


Fig. 2.22. Shapes of constant energy surfaces in the conductive zones of german and silicon

The values of the components of the tensor of the effective mass of an electron determined by cyclotron resonance are $m_1 = m_2 = m_3 = 0.082m_0$; $m_s = m_l = 1.69m_0$, and in silicon $m_1 = m_2 = m_3 = 0.19m_0$; $m_s = m_l = 0.92m_0$. m_s/m_l ratio characterizes the anisotropy of isoenergetic surface properties. For Germanium, the ratio is 19.3, and for silicon 4.8.

The minimum distance between the bottom of the permeability zone and the top of the stress zone is called the bandwidth zone. In silicon and germanium, the

energy extremes of electrons and spikes lie at different points in the Brilllluen zone (see Figure 2.21). The prohibited zone of germanium with sufficient purity at room temperature and normal atmospheric pressure is 0.67 eV, that of silicon is 1.12 eV, that of gallium arsenide is 1.43 eV, and that of indium antimony is 0.18 eV.

As the strong bond approaches, the tensile zone of germanium and silicon is formed from the p-orbitals of the atom, and therefore their tensile zone contains three sub-zones, which can be conventionally $P_{3/2}^{3/2}, P_{3/2}^{1/2}, P_{1/2}^{1/2}$ designated as -zones. The maximum energy for all three zones is located in the center of the Brilllluen zone (see Fig. 2.21). In this case, both 3/2-zones are excitation at $k = 0$. For the 1/2-zone, the excitation is partially removed due to the spin-orbit interaction, as a result of which it is positioned below E_{so} .

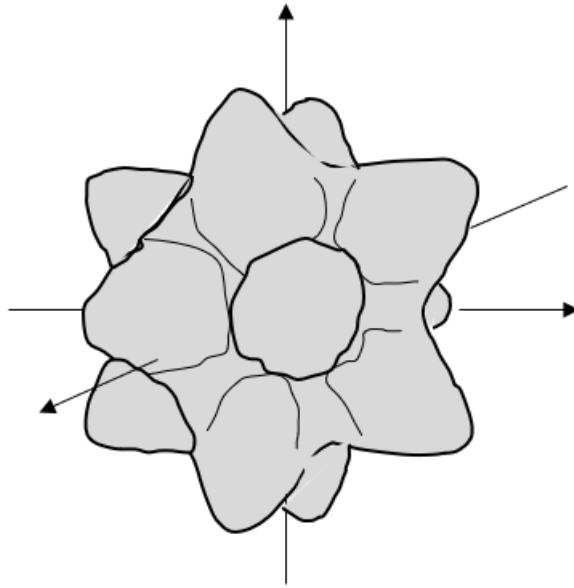


Fig. 2.23. Continuous Energy Surfaces of Heavy Swells in Silicon

For both $P_{3/2}$ -zones, $k \neq 0$ loses excitation at 0, and for them compute the dependence $E(k)$

$$E_{1,2}(k) = E(0) - \frac{\hbar^2}{2m_0} [Ak^2 \pm \sqrt{B^2k^4 + C^2(k_x^2k_y^2 + k_z^2k_y^2 + k_x^2k_z^2)}] \quad (2.153)$$

where A, B, C are variable magnitudes, equal to $13.1 \pm 0.4; 8.3 \pm 0.6$ for germanium; 12.5 ± 0.5 and $4.0 \pm 0.1; 1.1 \pm 0.4; 4.1 \pm 0.4$ for silicon. Surfaces of the same energy (2.153) represent its corrugated surfaces (Figs. 2.23 and 2.24).

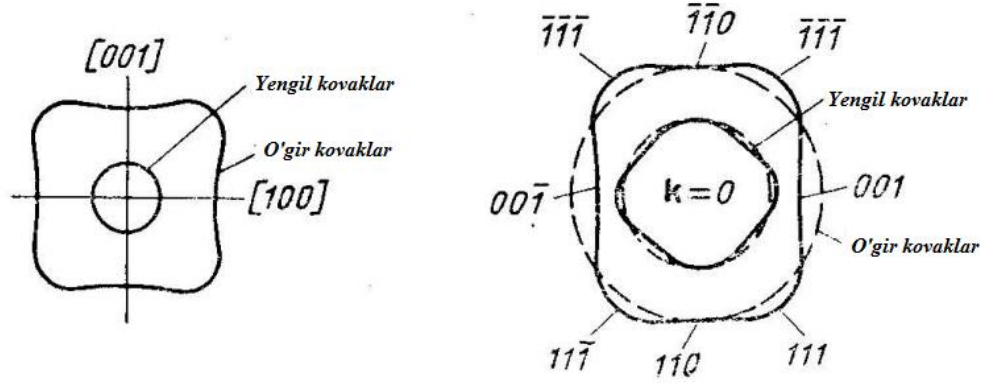


Fig. 2.24. Isoenergetic surfaces of germanium and silicon in valent zones

Averaging the dependence in different directions (2.153) in k -space allows the corrugated surface to be replaced by a sphere (Fig. 2.24), for which $E_{1,2}$

$$E_{1,2}(k) = E(0) - \frac{\hbar^2 k^2}{2m_0} \left[A \pm \sqrt{B^2 + \frac{1}{5} C^2} \right] \quad (2.154)$$

can be expressed in expression.

In this case, the effective mass will be the scalar magnitude

$$m_p^* = \frac{m_0}{A \pm \sqrt{B^2 + C^2 / 5}} \quad (2.155)$$

There are 2 types of hives

Og'ir Cabins

$$m_{p\tau}^* = \frac{m_0}{A - \sqrt{B^2 + C^2 / 5}} \quad (2.156)$$

Yengil kovaklar

$$m_{pa}^* = \frac{m_0}{A + \sqrt{B^2 + C^2 / 5}} \quad (2.157)$$

Based on Fig. 2.21, the zone of heavy hives will be wide ($\partial^2 E / \partial k^2$ will be small), while the zone of light hives will be narrow ($\partial^2 E / \partial k^2$ will be large).

For the effective masses of heavy and light canines from experiments on cyclotron resonance the following results were obtained:

For German

$$m_{p\tau}^* = 0,33m_0; \quad m_{pa}^* = 0,04m_0;$$

For silicon

$$m_{p\tau}^* = 0,56m_0; \quad m_{pa}^* = 0,16m_0;$$

For gallium arsenide

$$m_{p\tau}^* = 0,50m_0; \quad m_{pa}^* = 0,12m_0;$$

$E(k)$ dependency for the third branch of the valent zone

$$E_a(k) = E(0) - E_{s0} - \frac{\hbar^2 k^2}{2m_0} A \quad (2.158)$$

where E_{s0} - is the spin-orbital interaction energy, which is 0.3 for germanium and 0.04 for silicon. Based on (2.158), the equal-energy surface of the tertiary valent zone consists of a sphere. Therefore, the third type of effective mass of the hive is scalar magnitude and $m_p^* = m_0/A$

Experiments show that the permeability zone of indium antimonide has spherical symmetry, and its minimum is located at the center of the Brilluen zone (Fig. 2.25). Energy near the minimum obeys the parabolic law, and the curvature of the zone is very large. Therefore, the effective mass of electrons located at the bottom of this zone is very small and is $0,013m_0$ according to cyclotron resonance measurements. For high energies, the conductivity zone has a non-parabolic shape and its curvature decreases. This is confirmed by the data in Fig. 2.26, which shows that the energy is not proportional to k^2 . In this case, energy close to the minimum, the so-called nonstandard zone, looks like this:

$$E = E_c + \frac{\hbar^2 k^2}{2m_0^*} (1 - \alpha k^2) \quad (2.159)$$

where α is a parameter that determines the degree of deviation of the permeability zone from parabolic; E_c - energy of the bottom of the permeability zone; m_0^* is the effective mass of the electron.

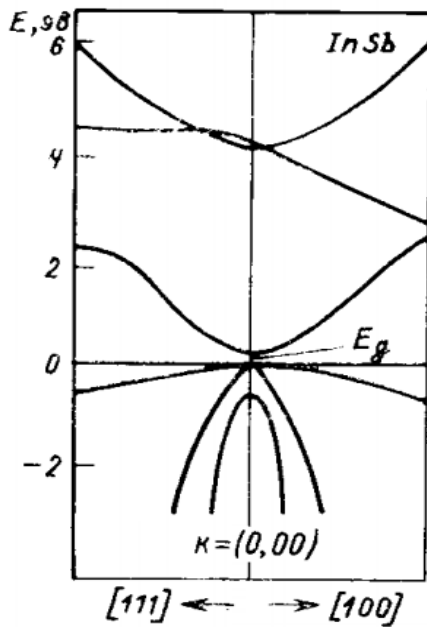


Fig. 2.25. Energy zone structure of indium antimony

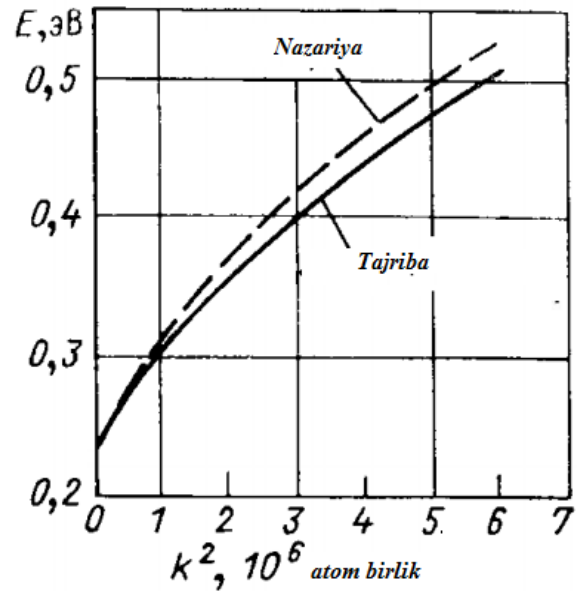


Fig. 2.26. Deviation of the permeability zone from the parabola in indium antimonide

The tensile zone of indium antimonide, the zonal structure of which is shown in Figure 2.25, contains three horns $E(k)$. The heavy hive zone V_1 ($m_{p\tau}^* = 0,6m_0$)

and the light hive zone V_2 ($m_{pa}^* = 0,012m_0$) were induced at $k = 0$. At $E_s = 0,9$ eV, the third zone is separated due to the spin-orbit interaction.

2.14. Effective mass method

When a semiconductor is alloyed with donors and capacitors, local states arise in it, which are characterized by additional layers. But before presenting the theory of localized states of an electron in a crystal, let us first consider a way to solve the Schrödinger's equation in an approximated way when an external bonding $U(r)$ is applied to the lattice area of the crystal $V(r)$. In this case, Schredinger's equation would be:

$$\left(-\frac{\hbar^2}{2m_0} \Delta + V(r) + U(r) \right) \Psi(r) = E \Psi(r) \quad (2.160)$$

This equation is difficult to solve directly, because the exact shape of the $V(r)$ field is not known.

From the previous consideration we know that if we apply the effective mass tensor m^* instead of the mass m_0 of the electron, then the motion of the electron will be the same as the motion of the classical free electron, but this electron is also affected by the external electric field in addition to the constant area of the crystal.

The Schredinger equation for a free electron is written as follows:

$$-\frac{\hbar^2}{2m_0} \Delta \Psi_0(r) = E \Psi_0(r) \quad (2.161)$$

The Schredinger equation for the electron in the crystal is written similarly to (2.161) as follows:

$$-\frac{\hbar^2}{2m^*} \Delta \Psi(r) = E \Psi(r) \quad (2.162)$$

Under what conditions (2.162) it is necessary to determine that Equation $V(r)$ gives the solution of the periodic field Schrödinger equation

$$\left(-\frac{\hbar^2}{2m^*} \Delta + V(r) \right) \Psi(r) = E \Psi(r) \quad (2.162)$$

As we know, the equation (2.163) Ψ is satisfied by the Blox functions $k(\mathbf{r}) = U k(\mathbf{r}) e i k r$, while the energy spectrum of the electron in the extreme is

$$E(k) = E(k_0) + \frac{\hbar^2}{2} \left[\frac{(k_x - k_{x0})^2}{m_1} + \frac{(k_y - k_{y0})^2}{m_2} + \frac{(k_z - k_{z0})^2}{m_3} \right] \quad (2.164)$$

is described by expression.

Equation (2.164) assumes that the inverse effective mass tensor is reduced to a diagonal form, so that Equation (2.162) can be written as follows

$$\sum_i \frac{1}{m_i} \frac{\partial^2 \Psi}{\partial x_i^2} + \frac{2E}{\hbar^2} \Psi = 0 \quad (2.165)$$

The solution to this equation

$$\Psi = Ae^{i(\alpha x + \beta y + \gamma z)} \quad (2.166)$$

it seems, here

$$E = \frac{\hbar^2}{2} \left(\frac{\alpha^2}{m_1} + \frac{\beta^2}{m_2} + \frac{\gamma^2}{m_3} \right) \quad (2.167)$$

Comparing the expressions (2.164) and (2.167), α assuming $\alpha = k_x - k_{0x}$, $\beta = k_y - k_{0y}$, $\gamma = k_z - k_{0z}$, and moving the head of the counting system of energy from point $k - k_0$ to point k_0 , the solutions of equations (2.162) and (2.163) give the same result for the energy spectra of the electron.

Thus, in determining the energy spectrum of an electron in a crystal in the presence of an external field

$$\left(-\frac{\hbar^2}{2m^*} \Delta + U(r) \right) \Psi(r) = E \Psi(r) \quad (2.168)$$

The equation can be used, where $V(r)$ has no unknown potential.

This method is called the effective mass method. The advantage of this method is that the components of an effective mass tensor at an extreme point can be determined experimentally, e.g. by cyclotron resonance experiments. The fact that the efficient mass method is only appropriate for electrons at extreme points of energy does not in most cases impose serious limitations on solving problems in semiconductor theory, since the concentrations of electrons (or spikes) are much smaller than the number of cases in the zone, so they will be located in the zones with minimal energy allowed.

2.15. Basic theory of additional cases

Until now, ideal crystals with a periodic structure have been considered. We will now consider semiconductors whose grids contain the simplest defects in the form of input atoms of group V or III of the periodic table.

Suppose one of the nodes in a silicon crystal has been replaced by a pentvalent arsenic atom. Four of them form covalent bonds with the nearest neighboring silicon atoms. The fifth electron, without being involved in covalent bond formation, interacts with a large number of silicon atoms, resulting in a weakly bonded with respect to the arsenic ion by Coulomb interaction forces. Its state is similar to that of an electron in a hydrogen atom. However, when solving the problem of the ionization potential of the input electron, two additional cases must be taken into account. First, this electron is located not only in the Coulomb field of the arsenic ion, but also in the periodic field of the crystal lattice. In this connection,

when describing its motion it is necessary to use the not real, but the effective mass m^* . Second, the Ze charge interaction of the electron and the positive arsenic ion occurs in a solid with relative dielectric conductivity ϵ , so the potential energy of the input electron must be included in the expression $U(r) \epsilon n$.

$$U(r) = -\frac{Ze^2}{4\pi\epsilon_r\epsilon_0 r} \quad (2.169)$$

where ϵ_0 is the electric constant.

(2.169) The relation is valid only for macroscopic point charges, i.e., if the distance between the interacting charges is greater than the crystal lattice constant, i.e., if the orbit of the input electron must cover a large number of crystal lattice nodes.

Given the above, the Schrödinger equation for the fifth electron of an arsenic atom that does not participate in covalent bonding in accordance with (2.161) takes the following form

$$\left(-\frac{\hbar^2}{2m^*} \Delta - \frac{Ze^2}{4\pi\epsilon_r\epsilon_0 r} \right) \Psi_a = E_n \Psi_a \quad (2.170)$$

Solving this equation in a similar way to the equation of the hydrogen atom, we get the following expression for the specific values of the electron energy of the donor mixture:

$$E_n = E_c - \frac{m^* Z^2 e^4}{9\hbar^2 \epsilon_r^2 \epsilon_0^2 n^2} \quad (2.171)$$

wherein the electron energy is calculated from the bottom of the conduction zone, n is the quantum number ($n \geq 1$), which determines the level of excitation at different levels of donor accession.

This formula can be rewritten as follows:

$$E_n = E_c - \frac{m_0 Z^2 e^4}{8\hbar^2 \epsilon_0^2} \left(\frac{m^*}{m_0} \right) \frac{1}{\epsilon_r^2 n^2} \quad (2.172)$$

If we put into this expression the values of e , m_0 , $\hbar$, and ϵ_0 , and express energy in electron-volts

$$E_n = E_c - \frac{13.52 Z^2}{\epsilon_r^2} \left(\frac{m^*}{m_0} \right) \frac{1}{n^2} = E_c - \frac{E_d}{n^2} \quad (2.173)$$

where 13.52- ionization energy of the hydrogen atom and

$$E_d = \frac{13.52 Z^2}{\epsilon_r^2} \left(\frac{m^*}{m_0} \right) \quad (2.174)$$

The basic state ($n=1$) energy of the ed- donor input atom.

It follows from (2.173) that the energy level of the donor input corresponding to the basic state ($n = 1$) will be located E_d below the bottom of the conductivity zone in the semiconductor's forbidden zone.

The ionization energy of the donor input is equal to E_d , (2.174) times less than the ionization energy of the hydrogen atom according to the equation. It also depends on the Z^2 . This means that the level of the doubly charged ion of the donor entry must be located in the prohibited zone below the level of the one-time charged ion of that access.

If $m^* = 0,25m_0$, e.g. in germanium, where $\varepsilon_r = 16$, the ionization energy of the atoms of the pentvalent mixture will be about 0.01 eV. The values of ionization energy obtained experimentally for donor mixtures are very much in line with the calculated value (Figure 2.27).

Application to the calculation of the orbit of the input electron of a hydrogen model

$$a = \frac{\varepsilon_0 h^2}{\pi e^2 m^*} \varepsilon_r n^2 = a_0 \left(\frac{m_0}{m^*} \right) \varepsilon_r n^2 \quad (2.175)$$

leads to the formula, where $a_0 = \varepsilon_0 h^2 / \pi e^2 m_0 = 0.053 \text{ nm}$ is the radius of the first Boron orbit of a hydrogen atom.

The radius of the first orbit of the donor entry electron in Germania is: $a_r = 64a_0 = 3,4 \text{ nm}$. If we consider that the lattice constant is 0.562 nm, this orbit encompasses about 200 lattice nodes. This ε_r confirms the possibility of considering the motion of an input electron in a medium with constant p .

Now let's look at the case of the accessor. Let's assume that one of the silicon lattice nodes has been replaced by a boron atom, which is an element of the third group of the periodic table. In this case, one electron is missing in the input atom to fill the covalent. Incomplete bonding - the cavity acts like a positively charged particle that provides electrical neutrality in this part of the crystal. Similar to the problem of the motion of an electron in the Coulomb field of a positive ion, the solution to the problem for this case gives the inherent values of energy.

$$E_p = E_v + \frac{m_0 Z^2 e^4}{8 h^2 \varepsilon_0^2} \left(\frac{m^*}{m_0} \right) \frac{1}{\varepsilon_r^2 n^2} = E_v + \frac{E_a}{n^2} \quad (2.176)$$

here

$$E_a = \frac{13.52 Z^2}{\varepsilon_r^2} \left(\frac{m^*}{m_0} \right) \quad (2.177)$$

i.e., the accessor input level is in the zone range above the ceiling of the voltage zone E_a at the value of E_v . According to calculations, for germanium it is 0.01 eV;

Thus, in covalent semiconductors such as germanium, the "small" layers of donor and recipient mixtures created by atoms of group V and III elements are very well characterized using the hydrogen model. The ionization energy of these mixtures is inversely proportional to ε_r^2 and proportional to the effective mass of the charge carriers.

Elements of groups I, II, VI, and VIII, which are introduced into the semiconductor as adjunct mixtures, form "deep" energy layers. These elements can enter the grid in the form of an atom (donor D or acceptor A), a single-charged ion [$D^{(1)}$ or $A^{(1)}$], and a double-charged ion [$D^{(2)}$] or [$A^{(2)}$]. In addition, the same input can generate both donor and acceptor layers.

Control Questions

1. Why is the Schrödinger equation necessary?
2. What does the Schredinger equation for crystal look like?
3. What is the meaning of an adiabatic approach?
4. What is a one-electron approximation?
5. Explain the convergence of strongly bonded electrons?
6. Kvaziimpul nima?
7. How is the electron bound to the wave vector at the bottom and top of the energy band?
8. How does an electron in a crystal move when exposed to an external electric field?
9. Effective mass nima?

Chapter III

OSCILLATION OF CRYSTAL LATTICE ATOMS

3.1. One-dimensional oscillations in a same-sex thread

As is known, atoms in a crystal do not move, they carry out heat oscillations in relation to their equilibrium state. These oscillations break the solid periodicity of the crystal structure. Also, these scatterings have a major impact on the propagation of charge carriers in the crystal structure and on many physical phenomena in semiconductors (discussed in Chapters 5 and 6).

The theory of the vibrations of atoms in a three-dimensional crystal is relatively complicated. Therefore, we first consider the vibrations of the atoms in a one-dimensional crystal and then generalize the obtained results to a three-dimensional crystal.

We first consider the propagation of longitudinal waves in a one-dimensional solid strand with a length density ρ . The x piece of the thread Δx is affected by the force F on the thread, where the displacement at the point x is $u(x)$ and the deformation resulting from this force

$$l(x) = \partial u / \partial x \quad (3.1)$$

deformation at the point $x + \Delta x$ of the thread

$$l(x + \Delta x) = l(x) + \partial^2 u / \partial x^2 \Delta x \quad (3.2)$$

If the force F causes a deformation, then the elastic modulus T is calculated as follows:

$$T = F / l.$$

Then the force acting on the x *elementary piece* of the thread

$$F \Delta x = T [l(x + \Delta x) - l(x)] = T \partial^2 u / \partial x^2 \Delta x \quad (3.3)$$

This force is equal to the acceleration of the mass of the elementary piece of the strand

$$F \Delta x = \rho \partial^2 u / \partial t^2 \Delta x \quad (3.4)$$

As a result, we get the equation of free oscillations of the thread:

$$\frac{\partial^2 u}{\partial t^2} = \frac{T}{\rho} \frac{\partial^2 u}{\partial x^2} \quad (3.5)$$

here

$$v = \sqrt{T / \rho} \quad (3.6)$$

The propagation speed of an elastic wave, however, it does not depend on the frequency.

(3.5) solve the wave equation,

$$u = A e^{i(qx - \omega t)} \quad (3.7)$$

We get the result, where A - is the amplitude of vibration, v is the vibrational frequency, ω -cyclic frequency, $q=2\pi/\lambda$ - is the number of λ waves, is the wavelength.

For an elastic wave propagating in a continuous strand, the frequency is proportional to the number of waves

$$\omega=vq \quad (3.8)$$

At the same time, the number of waves varies ∞ from 0 to 0, as a result of which the frequency also varies from 0 to 0 ∞ .

3.2. Oscillations in a uniatomic linear chain

We consider a linear chain of atoms with masses m , located at the same distance a . Suppose that the displacement of the n th atom is u_n . If the displacement of this atom from the equilibrium state is small, then the forces of interaction between atoms are treated as quasielastic forces, i.e., they will depend on the change in the distance between the atoms.

Therefore, only taking into account the interaction with closely adjacent atoms, the force acting on the n -atom is expressed as follows:

$$F_n = -\beta(u_{n+1} - u_n) - \beta(u_n - u_{n-1}) = \beta(u_{n+1} + u_{n-1} - 2u_n) \quad (3.9)$$

Here β is the coefficient of quasielastic force.

N -atom's chaotic equation

$$m d^2 u_n / dt^2 = \beta(u_{n+1} + u_{n-1} - 2u_n) \quad (3.10)$$

We look for the general solution of the equation in the following form:

$$u_n = A e^{i(qan - \omega t)} \quad (3.11)$$

By putting this solution into equation (3.10), we get:

$$-m\omega^2 = \beta(e^{iqa} + e^{-iqa} - 2) \quad (3.12)$$

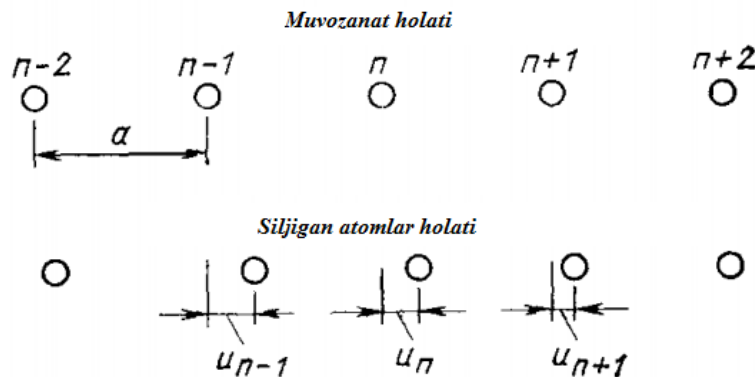


Fig. 3.1. A grid made of identical atoms

Given that $e^{iqa} + e^{-iqa} - 2 = -4\sin^2(aq/2)$

$$\omega = \pm \omega_m \sin(aq/2) \quad (3.13)$$

here

$$\omega_m = 2(\beta/m)^{1/2} \quad (3.14)$$

The bond between frequency and wave number - is called the law of dispersion. Hence, if the frequency of vibration of elastic waves is bound to the number of waves (3.13) according to the law of dispersion, the solution satisfies equation (3.10) for any n .

The propagation rate of an elastic wave along a linear chain depends on the frequency

$$v = \frac{\omega\lambda}{2\pi} = \frac{\lambda}{\pi}(\beta/m)^{1/2} \sin \frac{\pi a}{\lambda} \quad (3.15)$$

This is the inherent property of elastic waves in a discrete structured medium.

We define the limit of variation of the number of waves. To do this, we compare the wave numbers q and $q' = q + 2\pi g/a$, where g is integers.

$$(u_n)_g = A e^{i(qn - \omega t)} e^{i2\pi ng} = A e^{i(qn - \omega t)} = u_n \quad (3.16)$$

Since n - is the ordinal number of the atom, ng - is an integer $e^{i2\pi ng} = 1$.

From the expression (3.16) it follows that at all points and always $(u_n)_g = u_n$. This follows from the fact that the wave numbers q and q' are physically identical and describe a single vibration. Therefore, the dispersion relationship of the wave vector

$$-\pi/a \leq q \leq \pi/a \quad (3.17)$$

It is enough to see for the values.

From this it can be inferred that the domain where the wave vectors q lie coincides with the Brillouin zone for the wave vector k of electrons. Therefore, the values of q can be considered in the same space as k (or p).

To determine the value of q , we follow the steps. We arrange the atoms in a one-dimensional chain along the circle in such a way that the N - and 1-atoms are in equilibrium at the distance of the lattice constant a . The radius of the circle becomes larger because of the large number of atoms. Since only a few lattice constants in a chain interact, the conditions under which the boundary atoms are located do not affect the movement of the atoms within the chain. In this case, the boundary conditions can be replaced by the Born-Karman cyclic conditions, according to which

$$u_{n \pm N} = u_n, \quad (3.18)$$

because in this case $n \pm N$ - atom corresponds to the number of the n -atom.

(3.18) and (3.11) are just we impose:

$$u_{n \pm N} = A e^{i[qa(n \pm N) - \omega t]} = A e^{i(qn - \omega t)} e^{\pm iqaN} = u_n \quad (3.19)$$

From (3.19) it follows that $e^{\pm iqaN} = 1$. I mean

$$qaN = 2\pi g \quad (3.20)$$

from here

$$q = 2\pi g/aN \quad (3.20)$$

demak g

$$-N/1 \leq g \leq N/2 \quad (3.22)$$

It varies within boundaries.

Thus, for an atomic chain with N identical atomic, i.e., N degree of freedom, the wave vector q discretely changes and can assume a value $N\pi$ in the range from $-\pi/a$ to π/a , i.e., the greater the value of the elementary cells in the chain. Since N is large, the changes in q can be considered quasi-continuous.

The law of dispersion for a linear grid composed of identical atoms is shown in Fig. 3.2. The same figure shows the correlation $\omega(q)$ for elastic same-sex yarn.

For long waves, i.e., for the small values of q , the $\sin/2 \approx aq/2$ be considered.

In that case (3.13) and

$$\omega = 2(\beta/m)^{1/2} aq/2 = (\beta/m)^{1/2} aq \quad (3.23)$$

For a linear atomic chain, the linear density $\rho = m/a$, and the modulus of elasticity is equal to the ratio of the force to the relative change of displacement

$$T = \frac{F}{[u_n - u_{n+1}]/a} = \beta a \quad (3.24)$$

So that

$$v = \sqrt{T/\rho} = \sqrt{\beta/ma} \quad (3.25)$$

Taking into account (3.25) (3.23), the attitude takes the following form:

$$\omega = vq, \quad (3.26)$$

As for the maximum frequency

$$\omega_m = 2v\pi/a \quad (3.27)$$

Is equal to.

Thus, for long waves, the frequency of vibration of a single-atom chain is proportional to the number of waves q .

In wave dispersion, the propagating phase velocity of a monochromatic wave is $v_f = \omega/q$, and the group velocity of the wave packet is $v_g = d\omega/dt$ (Fig. 3.3). For a linear chain of identical atoms according to (3.13) and (3.27)

$$v_f = \omega_m \left| \frac{\sin aq/2}{q} \right| = v_t \left| \frac{\sin aq/2}{aq/2} \right| \quad (3.28)$$

$$v_g = \frac{a\omega_m}{2} |\cos aq/2| = v_t |\cos aq/2| \quad (3.29)$$

For long waves, i.e. for the smaller values of q , the phase and group velocities are equal and represent the velocity of sound

$$v_f = v_g = v_t \quad (3.30)$$

This means that at long wavelengths the uniatomic chain behaves like an elastic thread.

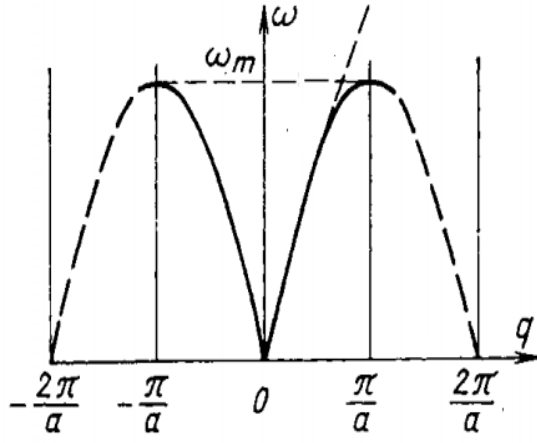


Fig. 3.2. For a linear grid, the frequency of vibration depends on the number of waves

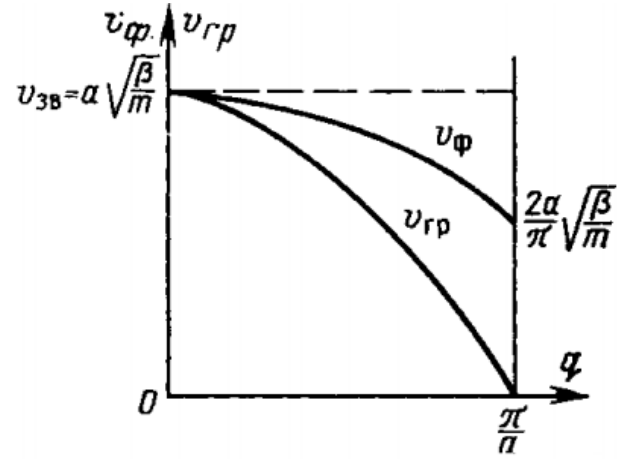


Fig. 3.3. The dependence of spatial and group velocity on the number of waves

In Fig. 3.2, as the values of q increase, i.e., with the decrease in wavelength, the frequency of oscillation of the uniaxial chain deviates from the linear dependence. The minimum wavelength corresponds to the maximum value of the number of wavelengths $q_m = \pi/a$ or $\lambda_{min} = 2a$. Physically, this case seems obvious, since in a discrete chain there would be no waves smaller than a half-wavelength lattice circuit. A wave with a minimum length $\lambda_{min} = 2a$ has a maximum frequency ω_m . Thus, the presence of the maximum frequency is inherent in the oscillations of discrete atomic structures. In solids, the sound velocity reaches $3 \cdot 10^5$ cm/s, and the lattice constant is about $3 \cdot 10^{-8}$ cm, so the maximum possible frequency of vibration in the lattice is $\omega_m = 2 \cdot 10^{13}$ Hz and lies in the infrared field.

For the shortest wavelengths, i.e., $q_m = \pi/a$, the phase velocity $v_f = 2v_t/\pi = 2a\sqrt{(\beta/m)}/\pi$ the group velocity (energy of lattice oscillations) decreases to zero. This shows that these oscillation modes describe stable waves in the grid.

3.3. The energy of vibration of atoms in a one-dimensional crystal. Normal coordinates

If 93.13) dispersion relation is performed, the expression (3.11) for a single-atom chain satisfies the action equation (3.10). But due to the strong binding forces of some atom, the vibrations are transmitted to neighboring atoms, and collective vibrations occur in the crystal that encompass all atoms. Therefore, in the case of small oscillations, the motion of the lattice atoms can be described as a linear superposition of running monochromatic waves. Each of these waves is

characterized by its own wave number q , frequency ω_q , and amplitude A_q . The displacement of each n th atom is as follows:

$$u_n = \sum A_q e^{i(qan - \omega t)} \quad (3.31)$$

Enter new coordinates

$$Q_q = C_q e^{-i\omega t} \quad (3.32)$$

where $C_q = \sqrt{BA_q}$. These coordinates change over time according to the harmonic law. These coordinates are called the normal (basic) coordinates of the crystal.

We rock the magnitudes of Q_q and A_q to (3.31)

$$u_n = \frac{1}{\sqrt{N}} \sum_q Q_q e^{iqan} \quad (3.33)$$

We determine the intrinsic (total) energy of the oscillations of the linear lattice atoms. It is equal to the sum of kinetic T and potentially U energies. Based on (3.33), kinetic energy is calculated as follows:

$$T = \sum_{n=1}^N \frac{m}{2} \dot{u}_n^2 = \frac{m}{2N} \sum_q Q_q^2 \sum_{n=1}^N e^{2iqan} \quad (3.34)$$

(3.21) Taking into account the expression $\sum_{n=1}^N e^{2iqan}$ We calculate the factor

$$\begin{aligned} \sum_{n=1}^N e^{2iqan} &= \sum_{n=1}^N e^{i \frac{4\pi}{N} ng} = e^{i \frac{4\pi}{N} g \cdot 1} + e^{i \frac{4\pi}{N} g \cdot 2} + \dots + e^{i \frac{4\pi}{N} g \cdot N} = \\ &= \frac{e^{i \frac{4\pi}{N} g} (1 - e^{i 4\pi g})}{1 - e^{i \frac{4\pi}{N} g}} e^{i \frac{4\pi}{N} g \cdot 1} = \begin{cases} 0, g \neq 0 \\ N, g = 0 \end{cases} \end{aligned} \quad (3.35)$$

Kinetic energy of one-dimensional lattice atoms given (3.35)

$$T = \sum_q \frac{m}{2} \dot{Q}_q^2 \quad (3.36)$$

Since $F_n = -\partial U / \partial u_n$, we get the following for the potential energy, taking into account (3.9) and (3.11)

$$U = U_0 + \frac{\beta}{2} \sum_{n=1}^N (2u_n - u_{n+1} - u_{n-1}) u_n \quad (3.37)$$

β To define (3.31), we put (3.10) into Equation (3.10)

$$m \sum_q \omega_q^2 A_q e^{i(qan - \omega_q t)} = \beta \sum_q (2 - e^{iqa} - e^{-iqa}) A_q e^{i(qan - \omega_q t)}$$

From here

$$\beta = \sum_q \frac{m \omega_q^2}{2 - e^{iqa} - e^{-iqa}} \quad (3.38)$$

Now we put (3.38) and (3.33) on (3.37)

$$U = U_0 + \sum_q \frac{m}{2} \omega_q^2 Q_q^2 \quad (3.39)$$

The Full Energy of the Grid

$$E = T + U = U_0 + \sum_q \frac{m}{2} (\dot{Q}_q^2 + \omega_q^2 Q_q^2) \quad (3.40)$$

When $U_0=0$

$$E = \sum_q \frac{m}{2} (\dot{Q}_q^2 + \omega_q^2 Q_q^2) = \sum_q E_q \quad (3.41)$$

origin.

$$E_q = \frac{m}{2} (\dot{Q}_q^2 + \omega_q^2 Q_q^2) \quad (3.42)$$

The expression corresponds to the expression of the total energy of the harmonic oscillator of mass m. Fixed energy values of a harmonic oscillator based on quantum mechanics

$$E_q = \hbar \omega_q (n + \frac{1}{2}), \quad n=0, 1, 2, \dots \quad (3.43)$$

Equality is determined by knowledge. Respectively, the lattice energy can be expressed in the form of the sum of the energies of non-dependent harmonic oscillators

$$E = \sum_q E_q = \sum_q \hbar \omega_q (n + \frac{1}{2}), \quad (3.44)$$

It follows from (3.41) that the total energy of the lattice oscillations, expressed in normal coordinates, has the appearance of the energy of a set of harmonic oscillators. This means that *a one-dimensional crystal with N atoms can be considered a dynamical system with degrees of freedom N and its motion can be described using normal coordinates independent of each other N normal coordinates. Each normal coordinate describes a specific configuration of all the atoms of a crystal oscillating according to a simple harmonic law (3.32). This collective vibration of crystal atoms is called the normal vibration of the lattice. In this, each atom of the crystal is involved in all N normal oscillations of the lattice. Q coordinates*

$$\ddot{Q}_q + \omega_q^2 Q_q = 0 \quad (3.45)$$

satisfies the equation, which is the equation of motion of the harmonic oscillator.

Thus, the total energy of the thermal motion of one-dimensional crystal atoms ω_q consists of the sum of the energies of the normal oscillations, which behave like linear harmonic oscillators.

3.4. Oscillations in a Two-Dimensional Chain of Two Atoms

Now, let's look at the oscillations in a system where there are two atoms in a single elemental grid of a crystal. Such materials may include ionic alkaline-halogen compounds, covalent semiconductors, germanium, silicon, as well as semiconductor compounds of $A^{III}B^V$ type.

Let's assume that a one-dimensional crystal consists of two types of atoms, one of which has a mass of m_1 and the other has a mass of m_2 (Fig. 3.4), atoms with a mass of m_1 are located at the n' lattice nodes, and atoms with a mass of m_2 are located at the n'' lattice nodes. We define the displacement of n' and n'' atoms with $u_{n'}$ and $u_{n''}$, and assume that the quasielastic forces between neighboring atoms are the same and β equal to a.

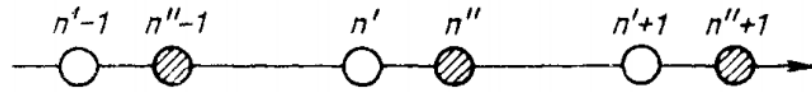


Fig. 3.4. A one-dimensional lattice with two atoms in an elementary cell

The equations of motion of the two types of atoms are:

$$\begin{cases} m_1 \ddot{u}_{n'} = \beta(u_{n''} + u_{n'-1}'' - 2u_{n'}) \\ m_2 \ddot{u}_{n''} = \beta(u_{n'} + u_{n-1}' - 2u_{n''}) \end{cases} \quad (3.46)$$

have views. Solutions of these equations

$$\begin{cases} u_{n'} = A_1 e^{i(qan - \omega t)} \\ u_{n''} = A_2 e^{i(qan - \omega t)} \end{cases} \quad (3.47)$$

So so. We put (3.47) on (3.46)

$$\begin{cases} (2\beta - m_1 \omega^2) A_1 = \beta(1 + e^{-iqa}) A_2 \\ (2\beta - m_2 \omega^2) A_2 = \beta(1 + e^{-iqa}) A_1 \end{cases} \quad (3.48)$$

It has a notrivial solution when the detegency of the system of equations A_1 and A_2 is zero

$$\begin{vmatrix} 2\beta - m_1 \omega^2 & \beta(1 + e^{-iqa}) \\ \beta(1 + e^{-iqa}) & 2\beta - m_2 \omega^2 \end{vmatrix} = 0 \quad (3.49)$$

$$e^{iqa} + e^{-iqa} = 2 \cos aq, \quad 1 - \cos aq = 2 \sin^2 aq / 2 \text{ For being}$$

$$\omega^4 - 2\beta \frac{m_1 + m_2}{m_1 m_2} \omega^2 + \frac{4\beta^2}{m_1 m_2} \sin^2 \frac{aq}{2} = 0 \quad (3.50)$$

we will have the equation. Solutions of this equation

$$\omega_1^2 = \frac{\omega_0^2}{2} \left[1 + \sqrt{1 - \gamma^2 \sin^2 \frac{aq}{2}} \right] \quad (3.51)$$

$$\omega_2^2 = \frac{\omega_0^2}{2} \left[1 - \sqrt{1 - \gamma^2 \sin^2 \frac{aq}{2}} \right] \quad (3.52)$$

here

$$\omega_0^2 = 2\beta \frac{m_1 + m_2}{m_1 m_2}; \quad \gamma^2 = 4 \frac{m_1 m_2}{(m_1 + m_2)^2} \quad (3.53)$$

From the formulas (3.51) and (3.52), the value of each q corresponds to the value of two ω . This means that in a one-dimensional crystal, when there are two atoms in its elementary cell, there are two types of oscillations: one is optical and the other is acoustic.

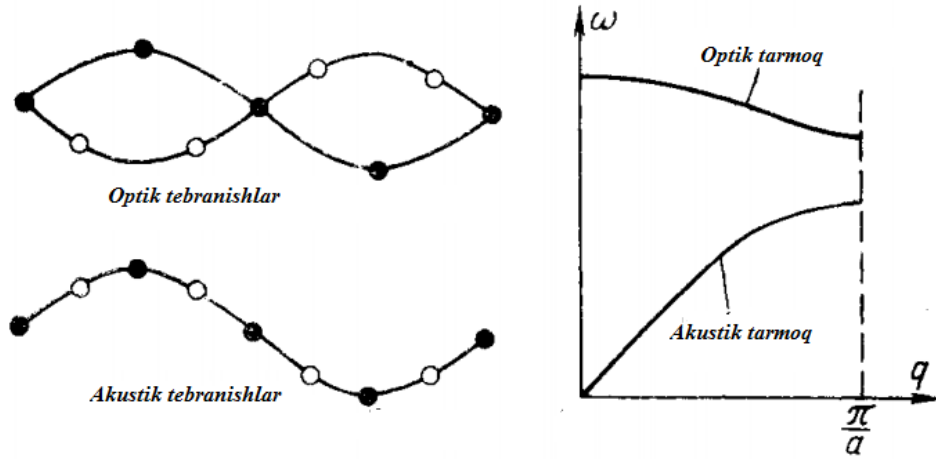


Fig. 3.5. Optical and acoustic vibrations with the same wavelength on a two-dimensional linear grid

We consider two extreme cases: the longest and the shortest wavelength. When $q=0$ and $q=\pi/a$

$$\begin{cases} \omega_{op}(0) = \omega_0; & \omega_{op}(\pi/a) = \frac{\omega_0}{\sqrt{2}} \sqrt{1 + \sqrt{-\gamma^2}} \\ \omega_{ak}(0) = 0; & \omega_{ak}(\pi/a) = \frac{\omega_0}{\sqrt{2}} \sqrt{1 - \sqrt{-\gamma^2}} \end{cases} \quad (3.54)$$

From here

$$\omega_{op}(0) = \omega_0 > \omega_{op}(\pi/a) > \omega_{ak}(\pi/a) > \omega_{ak}(0) = 0 \quad (3.55)$$

It turns out that there is a lot of work to be done.

For the longest wave of the acoustic network, when $\omega_{ak}(0)=0$, based on (3.47) and (3.48)

$$(u_n'/u_n'')_{ak} = (A_1/A_2)_{ak} = 1 \quad (3.56)$$

that is, atoms move the same distance.

This means that in an infinitely long wave of the acoustic network of oscillations, atoms move synchronously, and at any given moment in time their

deviations from the equilibrium state are the same, i.e., the cells move as a whole (Fig. 3.5).

A similar situation occurs in the propagation of elastic waves. That is why this network of vibrations was called acoustic.

The infinitely long wavelength of the optical oscillation network ω_0 coincides with a maximum frequency equal to. To do this, according to (3.47) and (3.48)

$$(u_n'/u_n'')_{op} = (A_1/A_2)_{op} = -m_2/m_1 \quad (3.57)$$

The response is justified.

Consequently, in an infinitely long wave of an optical network of oscillations, cell atoms move in the opposite direction, they oscillate in antiphase in such a way that the center of gravity of each pair of atoms remains motionless, i.e., $m_1 u_n' + m_2 u_n'' = 0$ (Fig. 3.5).

If a unit cell is composed of ions of different characters, then their oscillation in opposite phase causes the ions to shift and cause a change in the electric dipole moment of the cell. This would manifest itself in absorbing and emitting infrared radiation. That is why this type of oscillation is called optical.

For long waves, i.e., for the sub-values of q , it is possible to spread the roots in (3.51) and (3.52) into a series since $\sin(aq/2) \approx aq/2$. The result, for long wavelengths of the optical network of oscillations

$$\omega_{op} = \omega_0 \left(1 - \frac{\gamma^2 a^2}{32} q^2\right) \quad (3.58)$$

We will have it. $q \rightarrow 0$ The phase velocity of the optical oscillations tends to infinity, while the group velocity is zero.

For long wavelengths of acoustic network of oscillations

$$\omega_{ak} = a \sqrt{\frac{\beta}{2(m_1 + m_2)}} q = v_t q \quad (3.59)$$

The frequency of vibration is proportional to the sound velocity, phase and group velocities

$$v_f = v_g = v_t \quad (3.60)$$

equal to the speed of sound.

This means that the acoustic network of oscillations of lattice atoms at small values of q corresponds to ordinary sound waves, and their velocity of propagation is equal to the velocity of sound in a crystal.

For the shortest wavelength $\lambda = 2a$, i.e., $q = \pi/a$, the frequency of optical oscillations will be minimal:

$$\omega_{op}(\pi/a) = (2\beta/m_2)^{1/2} = \omega_{min} \quad (3.61)$$

And the frequency of acoustic vibrations is maximum:

$$\omega_{ak}(\pi/a) = (2\beta/m_2)^{1/2} = \omega_{max} \quad (3.61)$$

The relationship $\omega(q)$ for the optical and acoustic networks of one-dimensional lattice oscillations with two atoms in an elementary cell is shown in Figure 3.6. $\omega_{op} > \omega_{ak}$. As Q increases, the frequency of optical vibrations decreases slightly. When $q=0$ ω_{op} and ω_{ak} make the biggest difference.

3.5. Optical and acoustic vibrations

The established laws of vibration of a chain of one-dimensional atoms allow us to express general ideas about the thermal motion of atoms in a three-dimensional crystal. The displacement of atoms in a three-dimensional lattice can also be described by the superposition of the normal oscillations of the lattice. Every normal oscillation is a wave, and its direction of propagation is determined by the wave vector q . It is characterized by λ wavelength and frequency ω . The wave vector q has fixed values as the elementary cells in the crystal. In this case, *the values* q

$$-\pi \leq q_i a_i \leq \pi \quad (3.63)$$

It varies within boundaries. Here a_i is the grid constant ($i - x, y, z$).

A volumetric crystal $N = N_x N_y N_z$ contains elementary cells, each of which can contain an s atom. Since each atom has three levels of freedom, the crystal is characterized by a $3sN$ degree of freedom, and in such a crystal there will be waves of $3sN$ type.

Two types of elastic waves occur in a three-dimensional crystal; Longitudinal - waves of compression and thinning, and transverse - waves of shear deformation. For the first round, the direction of wave propagation is the same as the direction of vibration of the crystal atoms. Transverse waves propagating in a direction perpendicular to the plane of vibration of atoms have two main polar directions, since a transverse wave can be split into two independent transverse waves, in which the oscillations occur in two mutually perpendicular directions.

For a volumetric crystal containing an atom s in an elementary cell, there must be type $2s$ of transverse waves and type s of longitudinal waves, a total of $3s$ types of waves, i.e. $3s$ different vibrational networks. In other words, each value of q corresponds to $3s$ oscillations that differ in frequency and polarity. The three strands of vibrations are acoustic. For them, in the long-wave region, $v_f = v_g = v_t$ is appropriate. Hereby the frequency of longitudinal oscillations will be greater than the frequency of transverse oscillations, since the speed of longitudinal elastic waves is greater than the speed of transverse waves. The remaining $(3s-3)$ networks are optical, for which $q \rightarrow 0$ is $v_f \rightarrow \infty$ is, $v_g \rightarrow 0$. In general, for waves with varying polarizations, dispersion curves are different. At the same time, for an isotropic

crystal, the two transverse waves correspond to the same frequency values ω_q . Various horn oscillation excitations can also occur in anisotropic crystals for directions characterized by high degree of symmetry.

Fig. 3.7 shows the measured vibration spectra of the germanium, silicon, and gallium arsenide lattices. They have two atoms in their elementary cell, so these semiconductors have optical and acoustic horns of vibrations, and each of them has one longitudinal and two transverse horns of vibrations, which are spinning.

From Fig. 3.7, the frequency of acoustic vibrations tends to zero at $q \rightarrow 0$, which is typical for an acoustic wave. This is permissible for all three directions of oscillations - one longitudinal LA and two transverse TA (swing). For the small values of the wave vector q , the frequency of these oscillations is proportional to q for each direction of wave propagation, and the phase velocity is constant and has two different values for LA and TA. This is also characteristic of the acoustic waves in the crystal.

For all three directions of oscillations of the optical network, the frequency is maximum when $q \rightarrow 0$ and reaches 10^{13} Hz. $q \rightarrow 0$ the phase velocity tends to infinity and the group velocity tends to zero, which is characteristic of the optical network of oscillations. With an increase in q , the frequency of optical vibrations of the lattice decreases negligibly, and therefore it can be assumed that $\omega_{op} \approx \omega_0$, while the frequency of acoustic vibrations increases sharply. $q = \pi/a$ in the field, i.e., near the boundary of the Brillouin zone, the difference between optical and acoustic vibrations disappears.

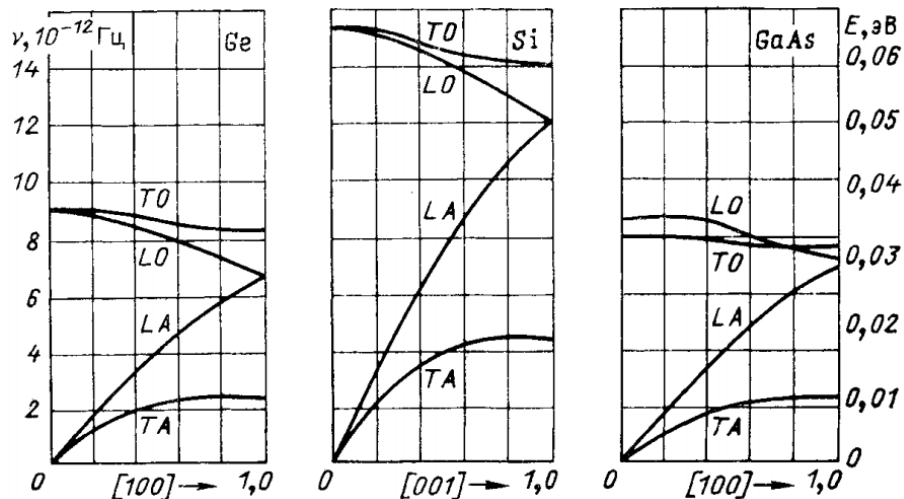


Fig. 3.7. Vibration spectra of germanium, silicon and gallium arsenide: TA and LA-longitudinal and transverse axial oscillations; TO and LO- TA and LA-longitudinal and transverse optical oscillations

Thus, the oscillation of the crystal atoms can be expressed as a superposition of the normal oscillations of $3sN$. But every normal oscillation, from a mechanical

point of view, is a harmonic oscillator. Consequently, the total energy of the crystalline oscillations is equal to the sum of the vibrational energies of the $3sN$ harmonic oscillators that do not interact with each other. From the point of view of quantum mechanics $\omega_j(q)$, energy is equivalent to an oscillator q_j oscillating with a frequency $E_{qj} = \hbar\omega_j(q)[n_{qj} + 1/2]$.

Thus, the total energy of the crystal oscillations

$$E = \sum_{qj} E_{qj} = \sum_{qj} \hbar\omega_j(q)[n_{qj} + 1/2], \quad (n_{qj}=0, 1, 2, \dots) \quad (3.64)$$

is defined by expression. where q is a wave vector of different value N ; j - Oscillation network (*there are 3s oscillation networks or 3s types of oscillations*).

$E_{qj} = \hbar\omega_j(q)[n_{qj} + 1/2]$ A quantum oscillator with energy $\Delta E_{qj} = \hbar\omega_j(q)\Delta n_{qj}$ can convert its energy into value. In this case, according to quantum mechanics, the change in the quantum number of the oscillator must be equal to:

$$\Delta n_{qj} = \pm 1 \quad (3.65)$$

$\hbar\omega_j(q)$ The energy quantum is called lattice oscillations, or phonon energy quantum. If $\Delta n_{qj} = +1$ the lattice goes into a high energy state. Therefore, the transition $\Delta n_{qj} = +1$ can be called the process of absorption of the phonon by the lattice, and $\Delta n_{qj} = -1$ can be called the process of emission of the phonon by the lattice. The phonon is therefore the only quantum excitation of a given normal vibration.

Similar to a photon - an electromagnetic field quantum, a phonon can be seen as a carrier of the energy quantum of lattice oscillations, i.e., a quasarra. Thus, the lattice vibrations are energetically equivalent to the background gas. In this case, $\Delta n_{qj} = +1$ represents the creation of the phonon and $\Delta n_{qj} = -1$ means the destruction of the phonon.

3.6. Phonons statistics

To determine the average number of energetic phonons $\hbar\omega_j(q)(n_{qj})$ present at temperature T at given normal oscillation, one needs to know the mean energy of the oscillator. So

$$n_{qj} = \frac{\langle E_{qj} \rangle}{\hbar\omega_j(q)} \quad (3.66)$$

Let the harmonic oscillator in heat equilibrium at temperature T be in the s -quantum state with energy E_s . As confirmed in statistical physics, the probability of an oscillator being in this state is:

$$W_s = Ce^{-E_s/kT} \quad (3.67)$$

Coefficient C is determined from the condition of probability normalization

$$\sum_s W_s = 1 \quad (3.68)$$

From this

$$C = (\sum_s e^{-E_s/kT})^{-1} \quad (3.69)$$

origin.

The mean or equilibrium energy of an oscillator is equal to the sum of its energies multiplied by the corresponding probabilities

$$\langle E \rangle = \sum_{n=0}^{\infty} E_n W_n \quad (3.70)$$

For the equilibrium energy value of the harmonic oscillator given (3.68) and

$$\langle E \rangle = \frac{\sum_{n=0}^{\infty} E_n e^{-E_n/kT}}{\sum_{n=0}^{\infty} e^{-E_n/kT}} = -\frac{d}{d(1/kT)} \ln \sum_{n=0}^{\infty} e^{-E_n/kT} \quad (3.71)$$

We'll have an expression.

Considering that it is the harmonic oscillator energy $E_n = \hbar\omega(n + \frac{1}{2})$, we define the sum in (3.71)

$$\sum_{n=0}^{\infty} e^{-E_n/kT} = e^{-\hbar\omega/2kT} (1 + e^{-\hbar\omega/kT} + e^{-2\hbar\omega/kT} + \dots) = \frac{e^{-\hbar\omega/2kT}}{1 + e^{-\hbar\omega/kT}} \quad (3.72)$$

We rock (3.72) by (3.71) and differentiate by $1/kT$

$$\langle E \rangle = \frac{\hbar\omega}{2} + \frac{\hbar\omega e^{-\hbar\omega/kT}}{1 - e^{-\hbar\omega/kT}} = \frac{\hbar\omega}{2} + \frac{\hbar\omega}{e^{\hbar\omega/kT} - 1} \quad (3.73)$$

The temperature-dependent $\hbar\omega/2$ voltage is called the zero energy of the oscillator. Since only normal oscillations excited above the zero energy level are involved in the interaction with electrons, zero energy can be ignored. Thus, we assume that the average energy of a harmonic oscillator is:

$$\langle E \rangle = \frac{\hbar\omega}{e^{\hbar\omega/kT} - 1} \quad (3.74)$$

In that case

$$\langle n \rangle = \frac{\langle E \rangle}{\hbar\omega} = \frac{1}{e^{\hbar\omega/kT} - 1} \quad (3.75)$$

It follows that phonons obey the Bose–Einstein statistics, $\langle n \rangle$ is the average number of phonons in a given normal oscillation, i.e. the $\hbar\omega$ average number of photons with energy in thermal equilibrium (their concentration) in a single quantum state – h^3 with energy in a volumetric phase space cell $\hbar\omega$.

At elevated temperatures, $\gg kT \hbar\omega$, the oscillator will be in the high energy state and the number of phonons in a cell of phase space is:

$$\langle n \rangle \approx kT / \hbar\omega \quad (3.76)$$

Mean energy of the oscillator

$$\langle E \rangle = \langle n \rangle \hbar\omega = kT \quad (3.77)$$

That is, corresponds to classical statistics.

At low temperatures, kT is $< \hbar\omega$

$$\langle n \rangle = e^{-\hbar\omega / kT} \quad (3.78)$$

Mean energy of the oscillator

$$\langle E \rangle = \hbar\omega e^{-\hbar\omega / kT} \quad (3.79)$$

Thus, the number of phonons in a quantum state depends on temperature. In addition, when electrons interact with or with each other, phonons are set up and disappear. Hence phonons are called quasi-particles because they do not behave like a gas of ordinary particles (atoms, electrons), but are independent of volume, temperature and number of particles.

In a solid, acoustic and optical vibrations are possible, but each lattice vibration is energy-equivalent to a phonon. Therefore, we can speak about acoustic and optical phonons. According to their mean number in the quantum state (3.75), the phonon energy is dependent on $\hbar\omega_j(q)$ and the temperature T .

$\omega_{op} > \omega_{ac}$ in the entire range of $\omega_{op} > \omega_{ac}$, so the energy of the optical phonons is greater than the energy of the acoustic phonons (see Fig. 3.7). Energy of acoustic phonons for small values of Q

$$E_{akf} = \hbar\omega_{jak} = \hbar\omega_{jtq}$$

The optical phonons are much smaller than their energy, for which $E_{opf} = \hbar\omega_{oj}$ is the maximum.

At low temperatures, there are no optical phonons when $< kT \hbar\omega_{oj}$, only acoustic phonons are excited and their number in one quantum state is determined by parity (3.78). At high temperatures, when $kT \gg \hbar\omega_{oj}$, in addition to acoustic phonones, a large number of optical phonons are induced.

3.7. Thermal capacity of the crystal grid

According to the definition, the thermal capacity of a solid

$$C = \partial E / \partial T \quad (3.80)$$

Since the harmonic oscillators do not interact, the total energy of the crystal at thermodynamic equilibrium is sum of the average values of the energy of all the oscillators. In this case, the sum must be made taking into account the types of vibrations that can be wave vector q and 3s of different values N . It is also necessary to consider zero energy as follows:

$$E_0 = \frac{1}{2} \sum_{jq} \hbar \omega_j(q)$$

Therefore, taking into account (3.74), the energy of lattice oscillations can be written as follows:

$$E = E_0 + \sum_{j=1}^{3s} \sum_q \frac{\hbar \omega_j(q)}{e^{\hbar \omega_j(q)/kT} - 1} = E_0 + \langle E_A \rangle + \langle E_O \rangle \quad (3.81)$$

Here is the balanced value of the energy of the acoustic vibrations of the lattice

$$\langle E_A \rangle = \sum_{j=1}^3 \sum_q \frac{\hbar \omega_j(q)}{e^{\hbar \omega_j(q)/kT} - 1} \quad (3.82)$$

And the equilibrium value of the energy of optical oscillations

$$\langle E_O \rangle = \sum_{j=1}^{3s} \sum_q \frac{\hbar \omega_j(q)}{e^{\hbar \omega_j(q)/kT} - 1} \quad (3.83)$$

We define $\langle E_A \rangle$ Since the wave vector q varies quascontinuously from 0 to π/a , and has a value of N (the number of elementary cells in a crystal), $\omega_j(q)$ also changes quascontinuously. In an acoustic network, it ranges from 0 to ω_{Amax} . Since N is so large, the sum of j in (3.82) can be substituted for integral

$$\langle E_A \rangle = \int_0^{\omega_{Amax}} dZ \frac{\hbar \omega d\omega}{e^{\hbar \omega/kT} - 1} \quad (3.84)$$

Here dZ - corresponds to the number of normal oscillations in the range from q to $q+dq$, $\omega_{and} Amax$ corresponds to q_{max} .

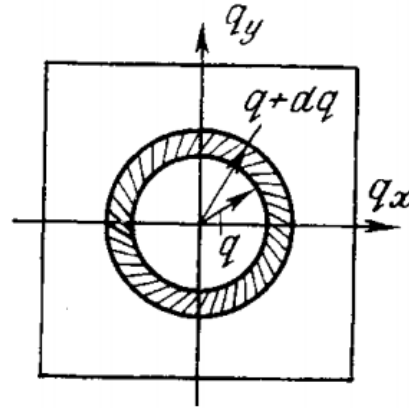


Fig. 3.8. The size of the layer between two layers

Three on the acoustic networks dZ To find the number of oscillations q -Radii in space q dan $q+dq$ We distinguish radial spherical surfaces up to 100. The size of this layer $dV_q = 4\pi q^2 dq$ (Fig. 3.8) According to (3.21), the size of the cubic crystal Bryulien zone elementary cell

$$\Delta V_B = dq_x dq_y dq_z = \frac{8\pi^3}{V} \quad (3.85)$$

where $V=a^3N_xN_yN_z=a^3N$ -crystal size, , $N= N_xN_yN_z$ is the number of elementary cells in a crystal.

Each Brulien zone elementary cell corresponds to one vibrational state of the crystal. To a single acoustic network of the interval dq respectively

$$dZ = V \frac{q^2 dq}{2\pi^2} \quad (3.86)$$

The amount corresponds normal vibration.

Based on Debay's theory, the law of dispersion would be appropriate for a triple acoustic network in the entire variable domain of q

$$\omega_j(q) = v_{jak} q \quad (j=1, 2, 3) \quad (3.87)$$

It is important to note that in this case the crystalline body is treated as an isotropic continuum in which there is no dispersion. Given (3.87), DQ is a single acoustic range in the interval, i.e., ω the number of oscillations in the range from frequency to $\omega+d\omega$

$$dZ = V \frac{\omega^2 d\omega}{2\pi^2 v_{jak}^3} \quad (3.88)$$

Now we determine the vibration state density of the grid, i.e., the number of normal oscillations corresponding to the unit frequency range per unit volume

$$g(\omega) = \frac{dZ}{V d\omega} = \frac{\omega^2}{2\pi^2 v_{jak}^3} \quad (3.89)$$

Number of holars in the range $\omega+d\omega$ for possible acoustic oscillations ω of three types (longitudinal with a velocity of one v_{lak} and transverse with a velocity of two v_{tak})

$$G(\omega) = \frac{\omega^2}{2\pi^2 v_{lak}^3} + \frac{2\omega^2}{2\pi^2 v_{tak}^3} = \frac{3\omega^2}{2\pi^2 v_{0ak}^3} \quad (3.90)$$

is defined by expression, where v_{0ak}

$$\frac{1}{v_{0ak}^3} = \frac{1}{3} \left[\frac{1}{v_{tak}^3} + \frac{2}{v_{tak}^3} \right] \quad (3.91)$$

is the average sound speed by directions and types.

In this case the energy of the acoustic vibrations for a V-bulk crystal

$$\langle E_A \rangle = \int_0^{\omega_{Amax}} dZ G(\omega) \frac{\hbar \omega d\omega}{e^{\hbar \omega / kT} - 1} = \frac{3V\hbar}{2\pi^2 v_{0ak}^3} \int_0^{\omega_{Amax}} dZ G(\omega) \frac{\hbar \omega d\omega}{e^{\hbar \omega / kT} - 1} \quad (3.92)$$

is defined through the integral.

$x = \hbar \omega / kT$, $\hbar \omega_{Amax} / k = \theta$ We define it. Then $x_{max} = \theta / T$,

$$\langle E_A \rangle = \frac{3V\hbar}{2\pi^2 v_{0ak}^3} \left(\frac{kT}{\hbar} \right)^4 \int_0^{\theta/T} \frac{x^3 dx}{e^x - 1} \quad (3.94)$$

The maximum frequency can be determined by the condition that the number of full acoustic vibrations is equal to $3N$

$$\langle \int_0^{\omega_{Amax}} dZG(\omega)d\omega = \frac{3V}{2\pi^2 v_{0ak}^3} \int_0^{\omega_{Amax}} \omega^2 d\omega = \frac{V\omega_{Amax}^3}{2\pi^2 v_{0ak}^3} = 3N \quad (3.95)$$

From here

$$\omega_{Amax} = v_{0ak} \left(\frac{6\pi^2 N}{V} \right)^{1/3} \quad (3.96)$$

$$\theta = \frac{\hbar\omega_{Amax}}{k} = \frac{\hbar v_{0ak}}{k} \left(\frac{6\pi^2 N}{V} \right)^{1/3} \quad (3.97)$$

magnitude – called the characteristic temperature of a solid, or Debay temperature is the acoustic oscillation energy of phonons where $k\theta = q_{max} = \pi/a$ For most solids, the Debay temperature lies in the range of 100-300 K.

given (3.97) and (3.94)

$$\frac{3V\hbar}{2\pi^2 v_{0ak}^3} \left(\frac{kT}{\hbar} \right)^4 = 9kTN \left(\frac{T}{\theta} \right)^3 \quad (3.98)$$

can be made to look like it. So that

$$\langle E_A \rangle = \frac{9NkT}{(\theta/T)^3} \int_0^{\theta/T} \frac{x^3 dx}{e^x - 1} = 3NkTD \left(\frac{\theta}{T} \right) \quad (3.99)$$

Debay function here

$$D \left(\frac{\theta}{T} \right) = \frac{3}{(\theta/T)^3} \int_0^{\theta/T} \frac{x^3 dx}{e^x - 1} \quad (3.100)$$

is defined through the integral.

Now let's define $\langle E_O \rangle$. Frequency of optical oscillations wave vector q changes little over the entire range of change, and therefore can be considered equal to their marginal value $\omega_j(q) = \omega_{0j}$. We enter the debay temperature corresponding to the boundary frequency of a given optical network

$$\theta_{0j} = \frac{\hbar\omega_{0j}}{k} \quad (3.101)$$

For most substances, θ_{0j} lies in the range of 100-800 K.

Using Fig. 3.101, we estimate the energy of optical vibrations

$$\langle E_O \rangle = \sum_{j=1}^{3s} \sum_q \frac{\hbar\omega_j(q)}{e^{\hbar\omega_j(q)/kT} - 1} = NkT \sum_{j=1}^{3s} \frac{\theta_{0j}/T}{e^{\theta_{0j}/T} - 1} \quad (3.102)$$

It is possible to illustrate in the figure because for each optical network the sum of q is equivalent to multiplying by N .

Hence the energy of the grid oscillations

$$E = E_0 + \langle E_A \rangle + \langle E_O \rangle = E_0 + NkT \left(3D\left(\frac{\theta}{T}\right) + \sum_{j=4}^{3s} \frac{\theta_{0j}/T}{e^{\theta_{0j}/T} - 1} \right) \quad (3.103)$$

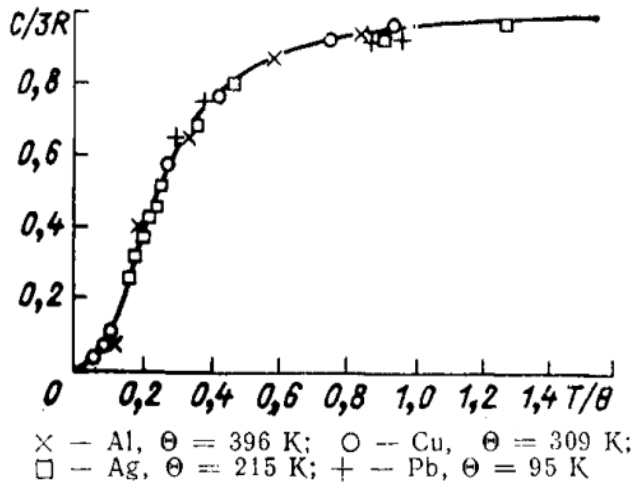


Fig. 3.9. Heat Capacity of Some Metals

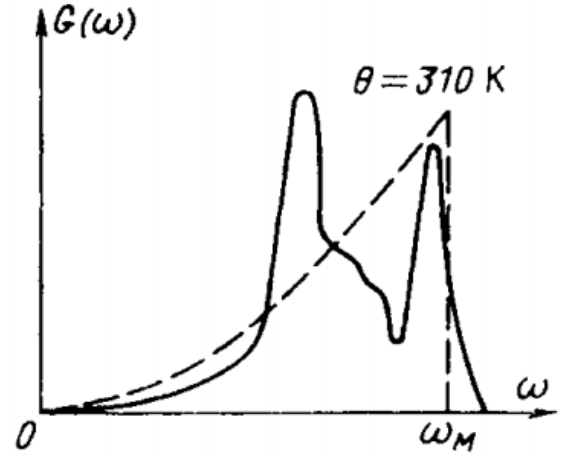


Fig. 3.10. Vibration density and Debay curve for tungsten grid (point)

In general, it is rather difficult to get an expression for heat capacity. We consider the sphere of high and low temperatures. $T \gg \theta_{0j}$, $T \gg \theta$ in the range of high and low temperatures

$$\begin{aligned} \text{ex} \approx 1+x, \quad 3D\left(\frac{\theta}{T}\right) &= 9\left(\frac{T}{\theta}\right)^3 \int_0^{\theta/T} x^2 dx = 3 \\ \sum_{j=4}^{3s} \frac{\theta_{0j}/T}{e^{\theta_{0j}/T} - 1} &= \sum_{j=4}^{3s} 1 = 3s - 3 \end{aligned} \quad (3.104)$$

So that

$$E = E_0 + 3RT \quad (3.105)$$

Where $R = 8.31 \text{ J/(mol} \cdot \text{K)}$ -Universal gas constant.

Heat capacity of the grille

$$C = \partial E / \partial T = 3R \quad (3.106)$$

This means that at high temperatures, the heat capacity of a solid does not depend on temperature.

When $T < \theta_{0j}$ is present at low tempears, the optical vibrations do not wake up and do not contribute to the energy of the lattice vibrations. When $T \ll \theta$ is (3.100), it can be put at the upper limit of the integral ∞ . Without it

$$\int_0^{\infty} \frac{x^3 dx}{e^x - 1} = \frac{\pi^4}{15} \quad (3.108)$$

Full charged energy,

$$E = E_0 + \frac{3\pi^4 NkT^4}{5\theta^3} \quad (3.109)$$

Heat capacity of the grid since it is equal to

$$C = \frac{12\pi^4 Nk}{5} \left(\frac{T}{\theta}\right)^3 \quad (3.110)$$

Thus, at low temperatures, the heat capacity of a solid is proportional to T^3 . The law of $C \sim T^3$ derived from Debay's theory is satisfied only in a few levels of domain close to absolute zero (Fig. 3.9). The deviation of the actual heat capacity from the Debay value (3.110) is determined by the fact that the actual vibrational density for the acoustic horns in the crystal differs greatly from the dependence $G(\omega)$ (3.90). This case is shown in Fig. 3.10, where the function $G(\omega)$ calculated for the tungsten lattice (continuous curve) and the Debay function $G(\omega)$ for the Debay temperature of tungsten equal to 310 K are shown (point curve).

3.8. Heat Expansion and Heat Resistance of a Solid

As we know, in the equation (3.37), only the quadratic displacement, i.e., the harmonic condition, was taken into account when determining the potential energy in the equation. Suppose this displacement equals x , then the potential energy is $U = \beta x^2$.

According to Boltzman, the probability of an atom deviating from its equilibrium state by a factor of x is:

$$f(x) = Ae^{-U/kT} \quad (3.111)$$

In this case, the distance between the atoms, which carry out harmonic oscillations, does not change, since their

$$\bar{x} = \frac{\int_{-\infty}^{\infty} xf(x)dx}{\int_{-\infty}^{\infty} f(x)dx} = \frac{\int_{-\infty}^{\infty} xe^{-\beta x^2/kT} dx}{\int_{-\infty}^{\infty} e^{-\beta x^2/kT} dx} = 0 \quad (3.112)$$

The mean displacement that can be determined by the formula is zero (due to the inaccuracy of the function under the integral, the displacement of the fraction is zero).

True bodies experience expansion from heat. This can be explained when we consider the cubic possibilities of expanding potential energy with the displacement of atoms. Calculating the anharmonic states of the interaction forces of atoms in a solid is a very complex task. However, considering a simple model of two interacting atoms allows us to understand the reasons for the expansion from heat. Furthermore,

this model makes it possible to determine the correct order of magnitude of the coefficient of linear thermal expansion.

Suppose two atoms move from the equilibrium state on the lattice to x distance from each other.

$$F = -\frac{dU}{dx} = -\beta x + 3\gamma x^2 \quad (3.113)$$

By force and interact with potential energy.

$$U = \beta x^2 - \gamma x^3 \quad (3.114)$$

If we put (3.114) into equation (3.111), then it is possible to spread the exponent corresponding to the anharmonic adjective in small displacements to the array, i.e. $e^{\gamma x^3 / kT} = 1 + \gamma x^3 / kT$. Without it

$$f(x) = A e^{\beta x^2 / kT} (1 + \gamma x^3 / kT) \quad (3.115)$$

Average displacement of the atom from equilibrium state

$$\bar{x} = \frac{\int_{-\infty}^{\infty} e^{-\beta x^2 / kT} (x + \gamma x^4 / kT) dx}{\int_{-\infty}^{\infty} e^{-\beta x^2 / kT} (1 + \gamma x^3 / kT) dx} = \frac{-\frac{\gamma}{kT} \frac{3\sqrt{\pi}}{4} (kT / \beta)^{5/2}}{(\pi kT / \beta)^{1/2}} = \frac{3\gamma kT}{4\beta^2} \quad (3.116)$$

Here the first part of the part under the integrals is zero, since the second part of the part is an odd function.

Known as linear expansion from heat

$$l_t = l_0 (1 + \alpha t) \quad (3.117)$$

Determined by the formula. Hence the temperature coefficient of linear expansion α is determined as the extension to units of length and 1 °C:

$$\alpha = \bar{x} / aT = 3\gamma / 4a\beta^2 \quad (3.118)$$

Mrs. Yerda A- Parametric Prison.

Thus, the temperature coefficient of the linear expansion of a solid is a constant magnitude; it is proportional to the anharmonicity coefficient and is zero at $\gamma = 0$.

The anharmonicity of the lattice oscillations makes it possible to explain the heat resistance of a solid body as well. In the case of harmonic oscillations of the lattice, the normal coordinates become independent, that is, subject to the principle of linear superpositions, according to which elastic waves propagate independently in the crystal, without scattering from each other. Since each normal coordinate representing a harmonic oscillator coincides on a given set of phonons, in a harmonic approximation the phonons do not interact, that is, they do not collide with each other. In this case, the heat flux – the energy carried away by phonons – dissipates at sound velocity ($v_g = v_t$), so that the heat resistance equals zero. When anharmonic

conditions in potential energy are taken into account, i.e., tertiary phenomena in the displacement of atoms, it makes the normal coordinates of the lattice, and consequently the oscillators, anharmonic and interrelated. The anharmonicity of the oscillators causes the phonons to scatter in the phonons. In this case, three phonons are involved in the collision process, whereby two phonons become one, or a single phonon splits in two. The scattering of phonons in the phonons together with the emergence and disappearance of phonons results in the emergence of thermal resistance in a solid. With increasing temperature, the anharmonicity of the vibrations increases since the amplitude of the vibrations increases and hence so does the heat resistance.

Control Questions

1. Explain the vibrational equation of a one-dimensional thread?
2. What are the oscillations of a single-atom linear chain?
3. What are normal coordinates?
4. What do you mean by vibrations of two- and three-dimensional lattice atoms?
5. What statistics are phonons subject to?
6. What is the heat capacity of a crystal grid?
7. What does the thermal expansion of a solid depend on?
8. What do you mean by the heat resistance of a solid?

Chapter IV.

STATISTICS OF ELECTRONS AND HIVES IN SEMICONDUCTORS

4.1. Quantum state density

The most important task of statistical physics is to determine the number of particles whose energy is located in a given range. To solve it, you need to know the number of quantum states and the probability of particles being in this state. Consequently, to determine the concentration of charge carriers in a semiconductor, one needs to know the true number of states occupied by electrons and hives.

Let there be quantum states dZ in a unit of volume crystal in the energy range from E to $E + dE$ (taking into account spin). We define the state density, i.e., the number of states in the unit energy range for a unit of crystal volume, by $N(E)$. Then by definition

$$N(E) = dZ/dE \quad (4.1)$$

If the probability of an electron filling the state E with energy is equal to $f(E, T)$, then the number of electrons located in the dZ states dZ is dZ

$$dn = f(E, T) dZ = f(E, T) N(E) dE \quad (4.2)$$

will be equal to . Accordingly, the number of electrons in the range of possible energy E_1 to E_2 will be equal to:

$$n = \int_{E_1}^{E_2} f(E, T) N(E) dE \quad (4.3)$$

Let us find the expression of the density of quantum states in the state where the equal-energy surfaces of the conduction band and the stress zone are spherical. We determine the density of cases at the lower edge of the permeability zone. According to (2.139), the energy of the electrons at the bottom of the zone

$$E = E_c + \frac{p^2}{2m_n^*} \quad (4.4)$$

can be written in the form, where $E_c = E(p_0)$ the energy of the electron at the bottom of the conduction zone; m^* is the effective mass of an electron.

We select a spherical layer wrapped between two isoenergetic surfaces corresponding to the energy of $E(p) = \text{const}$ and $E(p) + dE = \text{const}$ (Figure 4.1). The size of this layer

$$dV_p = 4\pi p^2 dp \quad (4.5)$$

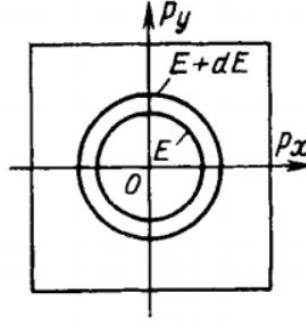


Fig. 4.1. The size of the Brillouin zone between the isoenergetic surfaces E and $E+dE$

Based on (2.79), the unit cell size of the Brillouin zone of the volume crystal in p -space is h^3 . Each cell can contain two electrons with an oppositely oriented spin. Given this, the number of quantum states in the volume of dV_p is equal to:

$$dZ = 2 \frac{dV_p}{h^3} = \frac{8\pi}{h^3} p^2 dp \quad (4.6)$$

Based on (4.4),

$$p^2 = 2m_n^*(E - E_c), \quad (4.7)$$

expression, from here

$$p = (2m_n^*)^{1/2} (E - E_c)^{1/2},$$

$$dp = \frac{1}{2} (2m_n^*)^{1/2} (E - E_c)^{-1/2} dE \quad (4.8)$$

expressions come from. By putting (4.6), (4.7), and (4.8) in relation (4.1), we get the expression of the density of quantum states at the bottom of a permeability zone with spherical symmetry:

$$N(E) = 4\pi (2m_n^*/h^2)^{3/2} (E - E_c)^{1/2} \quad (4.9)$$

We determine the density of cases located near the ceiling of the stress zone. To do this, we use (2.139) and (2.144) for the pit energy:

$$E = E_0 - \frac{p^2}{2m_p^*} \quad (4.10)$$

where $E_n = E(p_0)$ is the energy of the hive at the ceiling of the tensile zone; m_p^* is the effective mass of the hives. From calculations like the foregoing, we get the following:

$$N(E) = 4\pi (2m_p^*/h^2)^{3/2} (E_l - E)^{1/2} \quad (4.11)$$

Thus, if the energy of the charge carriers is a quadratic function of the quasi-impulse, then the density of the states $N(E)$ has an energy dependence of $(E - E_c)^{1/2}$ or $(E_v - E)^{1/2}$. Consequently, the expressions (4.9) and (4.11) apply only to cases close to the energy extreme, i.e. at the bottom and ceiling of the zone.

Now let's calculate the density of the quantum states for the complex conduction zone state when the $E(p)$ dependence at the bottom of the zone is described by the relation of (2.132). Because for the bottom of the zone, $p_{ox} = p_{oy} = p_{oz} = 0$ and $E(p_0) = E_c$, is then

$$E = E_c + \frac{p_x^2}{2m_1} + \frac{p_y^2}{2m_2} + \frac{p_z^2}{2m_3} \quad (4.12)$$

where $1/m_1, 1/m_2, 1/m_3$ are the diagonal components of the inverse effective mass tensor.

In this case, the isoenergetic surfaces are ellipsoids, and their equation in canonical form has the following form:

$$\frac{p_x^2}{a^2} + \frac{p_y^2}{b^2} + \frac{p_z^2}{c^2} = 1 \quad (4.13)$$

Here are the half axes of the ellipsoid

$$\left. \begin{aligned} a &= [2m_1(E - E_c)]^{1/2}; \\ b &= [2m_2(E - E_c)]^{1/2}; \\ c &= [2m_3(E - E_c)]^{1/2} \end{aligned} \right\} \quad (4.14)$$

The size of an ellipsoid with half axes a, b, c is equal to:

$$V_p = \frac{4\pi}{3} abc = \frac{8\pi}{3} (2m_1 m_2 m_3)^{1/2} (E - E_c)^{3/2} \quad (4.15)$$

The size of the layer wrapped between the ellipsoids of equal energy $E = \text{const}$ and $E + dE = \text{const}$ is:

$$dV_p = 4\pi (2m_1 m_2 m_3)^{1/2} (E - E_c)^{1/2} dE. \quad (4.16)$$

The dV_p volume, taking into account the spin, includes the following number of quantum states:

$$dZ = 2 \frac{dV_p}{h^3} = 4\pi \left(\frac{2}{h^2} \right)^{3/2} (m_1 m_2 m_3)^{1/2} (E - E_c)^{1/2} dE \quad (4.17)$$

Therefore, the state density expression at the bottom of the complex conductivity zone takes the following form:

$$N(E) = 4\pi (2/h^2)^{3/2} (m_1 m_2 m_3)^{1/2} (E - E_c)^{1/2} \quad (4.18)$$

If

$$m_1 m_2 m_3 = m_{dn}^{*3}, \quad (4.19)$$

Considering that this expression can be rendered in a form similar to the expression (4.9), where M_{DN}^* is called the effective mass of the state density for electrons.

In this case, for the ordinary zone

$$N(E) = 4\pi (2m_{dn}^*/h^2)^{3/2} (E - E_c)^{1/2} \quad (4.20)$$

Thus, if we introduce the effective mass of the state density for electrons computed by (4.19), the ellipsoid surface of the conduction zone of equal energy can be replaced by a sphere.

We determine the effective masses of the state density for electrons and spikes in silicon and germanium.

Silicon is located at a distance of $k = 0$ points in the direction of the absolute minimum energy in the conductivity zone of silicon from Fig. 2.21 [100], which is approximately 75% of the value of $2\pi/a$. In total, the first Brillouin zone contains six ellipsoids of equal energy stretching along the main crystallographic axes for an energy slightly greater than the E_c value (see Fig. 2.22). Accordingly, the density of cases expressed by equation (4.18) for an ellipsoid increases 6-fold. Given that for silicon $m_1 = m_2$ then

$$N(E) = 6 \cdot 4\pi(2/h^2)(m_1^2 m_3)^{1/2} (E - E_c)^{1/2} \quad (4.21)$$

and $m_1 = 0,19m_0$ va $m_3 = 0,92m_0$ effective mass of state density for electrons given values and

$$m_{dn}^* = 6^{2/3}(m_1^2 m_3)^{1/3} = 1,08m_0 \quad (4.22)$$

Consequently, all six ellipsoids of the isoenergetic surface of the conduction zone in silicon can be replaced by a single spherical surface with an effective mass of state density of $1,08m_0$ for electrons.

At point $k=0$, the tensile zone of silicon is broken twice, and the equal-energy surfaces have a much more complex shape, which is characterized by the formula (2.153). At the same time, they can converge with two spherical surfaces of the form, corresponding to the effective masses of heavy and light canines (2.154).

In this case, the density of cases is determined by the sum of the density of cases from heavy and light canine zones:

$$N(E) = 4\pi(2/h^2)^{3/2} (m_{pT}^{*3/2} + m_{p\pi}^{*3/2}) (E_v - E)^{1/2} \quad (4.23)$$

The isoenergetic surfaces of both lines can be replaced by a single reduced sphere with state density.

$$N(E) = 4\pi(2m_{dp}^*/h^2)^{3/2} (E_v - E)^{1/2} \quad (4.24)$$

To do this, the effective mass of case density for hives is equal to:

$$m_{dp}^* = (m_{pT}^{*3/2} + m_{p,1}^{*3/2})^{2/3} = 0,56m_0 \quad (4.25)$$

Here $m_{pT}^* = 0,56m_0$, $m_{p,1} = 0,16m_0$.

Germany's first Brillouin zone has eight energy minimums located at points $[1/2, 1/2, 1/2]$. Corresponding surfaces of equal energy are shown in Figure 2.22, which shows eight energy ellipsoids cut in half by the edges of the Brillouin zone. The two opposite halves of the ellipsoids in Fig. 2.22 can be considered to be a whole ellipsoid surface. Thus, there are four complete ellipsoids in the main Brillouin zone of the germanium, for which $m_1 = m_2$. Thus, the density of the states at the top of the germanium permeability zone is determined by the expression (4.20) with the effective mass of the state density for electrons:

$$m_{dn}^* = (4^2 m_1^2 m_3)^{1/3} = 0,56 m_0 \quad (4.26)$$

Here $m_1 = 0,082 m_0$, $m_3 = 1,59 m_0$.

The structure of the germanium valent zone is similar to that of the silicon zone (see Fig. 2.21). The effective mass of state density for spurs describing the density of quantum states in the upper part of the germanium valent zone:

$$m_{dp}^* = \left(m_{pT}^{*3/2} + m_{p\pi}^{*3/2} \right)^{2/3} = 0,35 m_0 \quad (4.27)$$

because $m_{pT}^* = 0,33 m_0$; $m_{p\pi} = 0,04 m_0$.

4.2. Fermi-Dirac section

In order to determine the number of particles having energy in a given range, in addition to the density of quantum states $N(E)$, it is necessary to know the probability of a particle occupying a particular state where energy is E , i.e., to know the distribution function $f(E)$. For particles with semi-integer spin obeying the Pauli principle under thermal equilibrium conditions, the Fermi-Dirac distribution would be appropriate.

$$f_0(E) = \frac{1}{e^{(E-F)/kT} + 1}, \quad (4.28)$$

where k is the Boltzman constant; T is absolute temperature; F is the Fermi energy or electrochemical potential, i.e., the work that must be spent to change the number of particles in a system by one unit under constant volume and temperature conditions.

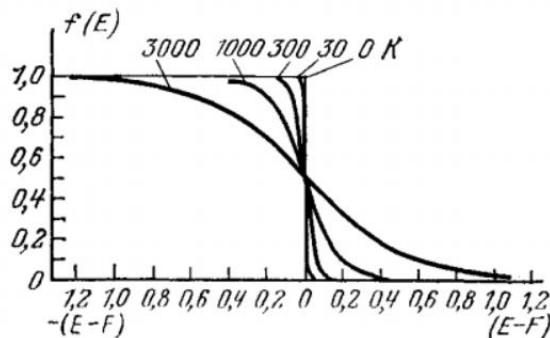


Fig. 4.2. Fermi–Dirac distribution

We consider the shape of the Fermi-Dirac distribution function at different temperatures. It follows from (4.28) that in the case of $T = 0$, in the energy range $0 < E < F$ is $f_0=1$ and $f_0=0$. This means that all quantum states with less energy than the Fermi level are occupied by electrons, and that the states above the Fermi level are completely free, not occupied by electrons. Fermi energy, then, is the maximum possible energy of electrons in a metal at absolute zero temperature.

Consider a case where $T > 0$. For an energy value equal to Fermi energy ($E = F$), we get $f_0 = 1/2$. Thus, the Fermi level is the energy level, and the probability of filling it at a temperature other than absolute zero is 0.5. At $T > 0$, a fraction of the electrons pass into states with a greater energy than Fermi energy ($E > F$) as a result of thermal motion, and accordingly a portion of the states below the Fermi level are free. In this case, the number of particles passing to high energy levels is equal to the number of free states generated in the $E < F$ region.

For the state $T > 0$, we estimate the change field of the distribution function $f_0(E)$. To do this, we calculate $f_0(E)$ for the different energy values. For energies varying from F to $\pm kT$, the value of $f_0(E)$ is $(1 + e)^{-1} = 0,27$, $(1 + 1/e)^{-1} = 0,73$. The f_0 values for $E - F = \pm 2kT$ are 0.118 and 0.882, and for $E - F = \pm 3kT$, the f_0 is 0.047 and 0.953. These data show that the probability of filling cases varies significantly from one or zero within $(2 \div 3)kT$ close to a value of $E = F$ (Figure 4.2).

The Fermi-Dirac distribution function describes the probability of filling a given quantum state with an electron. In the case where E has energy, the absence of an electron in the heat equilibrium, i.e., the probability that it is occupied by a hole will be equal to:

$$f_{0p}(E) = 1 - f_0(E) = 1 - \frac{1}{e^{(E-F)/kT} + 1} = \frac{1}{e^{(F-E)/kT} + 1} \quad (4.29)$$

Therefore, if we measure energy, the distribution function for the holes is similar to the distribution function for electrons. In the opposite direction to the direction of calculating the energy for electrons, the expression f_0 for electrons in cases where the energy of the particles is E from the Fermi level looks like this:

$$f_0(E) = e^{(F-E)/kT} = C e^{-E/kT}, \quad (4.30)$$

i.e. corresponds to the Boltzmann distribution function for particles subject to classical statistics. If charge carriers obey Boltzmann's statistics, then the electron gas is not degenerative, and accordingly, a semiconductor with such a distribution of charge carriers is generally called nondegenerative.

4.3. The level of replenishment of access levels

Consider a semiconductor containing a donor mixture with an N_d concentration. The donor holding the electron is electroneutral. This corresponds to, for example, the position where one of the locations of the silicon crystal lattice is occupied by an arsenic atom. In this case, the fifth tensile electron of the donor input atom is not involved in the covalent bonding and corresponds to the energy level located below the bottom of the conduction zone in terms of E_d value (Fig. 4.3-a).

Since the donor input has only one electron that can participate in the input conduction, the total number of cases for the donor input must be equal to the number of atoms of the mixture included in the unit volume of the crystal, i.e., N_d .

Suppose that the concentration of electrons on the donor input surface is n_d . In this case, p_d is the concentration of ionized donor atoms formed as a result of thermal transfers of electrons from the donor layer to the permeability zone and having a positive charge

$$p_d = N_d - n_d. \quad (4.31)$$

If, according to Pauli's principle, two electrons with an antiparallel spin can be placed on the donor input surface, then the probability of its filling is determined by the Fermi-Dirac function (4.28), in which it would be necessary to put E_d instead of E .

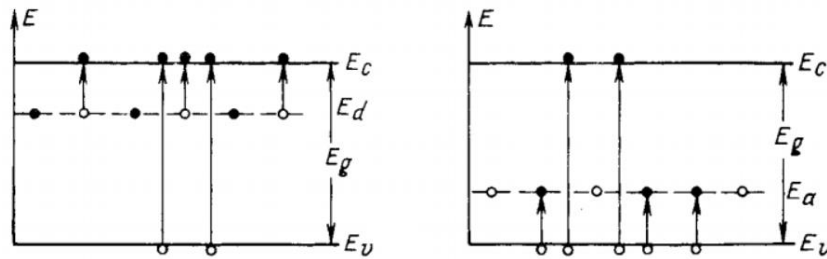


Fig. 4.3. Electronic (a) and receiver (b) semiconductors

E_d is the electron energy at the introductory level. But *there* can only be one electron in the E_d state, which can be captured in two ways, depending on the direction of rotation. Consequently, the neutral state of the donor input has twice the statistical weight of the ionized state. Since the probability of such a case is 1 if the donor does not have an electron at the input level, based on Boltzman's principle, we can write:

$$p_d : n_d = 1 : 2e^{-(E_d - F)/kT}. \quad (4.32)$$

Using (4.31) it is possible to write this parity in the form

$$\frac{N_d - n_d}{n_d} = (2e^{-(E_d - F)/kT})^{-1}, \quad (4.33)$$

It follows that the concentration of electrons on the donor entry surface is equal to:

$$n_d = \frac{N_d}{\frac{1}{2}e^{(E_d - F)/kT} + 1} \quad (4.34)$$

Concentration of positive ions of the donor mixture based on equations (4.32) and (4.34)

$$p_d = \frac{N_d}{2e^{(F-E_d)/kT} + 1} \quad (4.35)$$

In relation to

In the general case (4.34), the pre-exponential factor can be written as g^{-1} . Then the probability of finding an electron on the donor surface having energy E_d is determined by the expression, and the distribution function for the positive ions of the donor mixture based on (4.29) and (4.35) is as follows:

$$f_0(E_d) = \frac{1}{g^{-1}e^{(E_d-F)/kT} + 1} \quad (4.36)$$

From (4.36), it follows that the entry level for single-valent donor impurity is twice degenerated, the factor (level) of spin degeneration is $g = 2$.

Consider an accetor semiconductor, e.g. silicon fused with boron. Suppose the concentration of the input input is N_a . The energy diagram of such a semiconductor is shown in Fig. 4.3- b.

$$f_{0p} = \frac{1}{ge^{(F-E_d)/kT} + 1} \quad (4.37)$$

A neutral boron atom forms three covalent bonds with neighboring silicon atoms; An electron from a neighboring silicon atom can pass into this incomplete bond, and this requires energy equal to E_a . As a result, a free pit is formed, and the boron atom becomes a negatively charged boron ion. Thus, there is one electron with an optional spin direction at the energy level of the receiving mixture (neutral state of the accetor mixture) or two electrons with antiparallel spin (ionized state of the accetor mixture) when the atom of the accetor mixture captures an electron from the tensile zone to complete the pair bonding. Consequently, the degeneration rate of the accetor level is $g = 2$. Therefore, at a given temperature, the concentration of electrons at levels of receiver impurity is determined by the ratio of neither (or the concentration of negative ions) to the form.

$$n_a = \frac{N_a}{ge^{(E_a-F)/kT} + 1} \quad (4.38)$$

Concentration of hives at the aseptor inlet

$$p_a = \frac{N_a}{g^{-1}e^{(F-E_d)/kT} + 1} \quad (4.39)$$

4.4. Concentration of electrons and hives in zones

In an input semiconductor, for example, let us consider the process of the formation of free charge carriers under the influence of heat in a donor semiconductor, shown in Energy Diagram 4.3.

At absolute zero temperatures, electrons occupy states with minimal energy value. In other words, at $T = 0$ K, all states of donor input and the stress zone are busy with electrons, so there are no electrons in the conduction zone.

With an increase in temperature, electrons are separated from the donor inlet and the covalent bonds of the main substance are disrupted, resulting in free electrons and spurs. The free charge carriers, which occur during thermal excitation and are in thermodynamic equilibrium with a lattice, are called equilibrium or heat, and the process of their formation is called thermal formation (or ionization).

Consequently, in the case of thermal excitation in the donor semiconductor, the transition of electrons from the donor layer and the stress zone to the permeability zone occurs, where equilibrium hives are formed. These transitions are illustrated by the arrows in Figure 4.3. The inverse process of transitioning electrons from the conduction zone to the donor entry layer or stress zone is called recombination.

In an accuterial semiconductor, electrons and bubbles arise as a result of the thermal transition of electrons from the tensile zone to the accutor input layer and the conductivity zone (Figure 4.3-b).

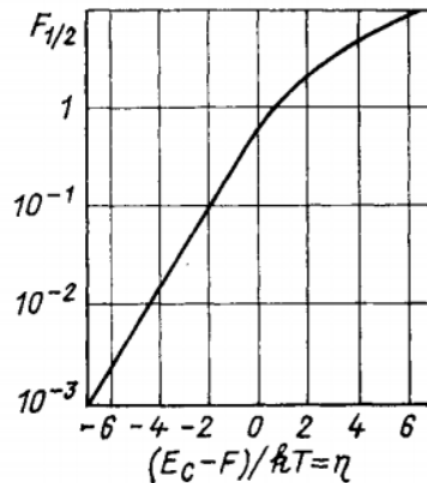


Fig. 4.4. Dependence of the Fermi–Dirac integral on the given Dirac dimension

For an input semiconductor, we determine the concentration of electrons in the n_0 conduction zone and the concentration of spurs in the stress zone p_0 in thermodynamic equilibrium. We solve the problem for a semiconductor having spherical surfaces of equal energy in both zones. If we replace the density of the quantum states in (4.3) with the state of isoenergetic surfaces (4.9) and use the Fermi-Dirac distribution function (4.28), the equilibrium electron concentration can be found. The integration must be done from the bottom of the E_c permeability zone to its top. If we consider that $f_0(E)$ is strongly dependent on energy, then the upper bound of the integral can be equal to ∞ . As a result, we get:

$$n_0 = \int_{E_c}^{\infty} f_0(E) N(E) dE = \int_{E_c}^{\infty} \frac{4\pi(2m_n^*/h^2)^{3/2}(E-E_c)^{1/2} dE}{e^{\frac{E-F}{kT}} + 1} \quad (4.40)$$

We enter dimensionless quantities

$$\frac{E-E_c}{kT} = \varepsilon; \quad \frac{F-E_c}{kT} = \eta \quad (4.41)$$

The magnitude is called the Fermi level. In this case (4.40) it looks like the following

$$n_0 = 4\pi \left(\frac{2m_n^*}{h^2} \right)^{3/2} (kT)^{3/2} \int_0^{\infty} \frac{e^{1/2} d\varepsilon}{e^{\varepsilon-\eta+1}} = N_c F_{1/2}(\eta) \quad (4.42)$$

$$N_c = 2 \left(\frac{2\pi m_n^* kT}{h^2} \right)^{3/2} \quad (4.43)$$

is the density of effective states in the permeability zone;

$$F_{1/2}(\eta) = \frac{2}{\sqrt{\pi}} \int_0^{\infty} \frac{\varepsilon^{1/2} d\varepsilon}{e^{\varepsilon-\eta+1}} \quad (4.44)$$

Fermi-Dirac integral. It is shown in Fig. 4.4.

The density of effective cases depends on the temperature N_c . If we put the infinitive values of the universal constants in (4.43),

$$N_c = 2(2\pi m_n^* kT/h^2)^{3/2} = 4,82 \cdot 10^{15} (m_n^*/m_0)^{3/2} T^{3/2} \quad (4.45)$$

We also find the concentration of hives in the valent zone. Based on (4.11) and (4.29)

$$p_0 = \int_{E_{vmin}}^{E_v} f_{0p}(E) \Lambda(E) dE = 4\pi \left(\frac{2m_p^*}{h^2} \right)^{3/2} \int_{-\infty}^{E_v} \frac{(E_v-E)^{1/2} dE}{e^{(F-E)/kT} + 1} \quad (4.46)$$

Here the lower bound of the integral has been replaced by ∞ because f_{op} is strongly dependent on energy. We enter a markup:

$$\frac{E_v-E}{kT} = \varepsilon_p; \quad \frac{E_c-E_v}{kT} = \frac{E_g}{kT} = \varepsilon_i; \quad \frac{F-E}{kT} = \eta + \varepsilon_i + \varepsilon_p \quad (4.47)$$

$$p_0 = \frac{4}{\sqrt{\pi}} \left(\frac{2\pi m_p^* kT}{h^2} \right)^{3/2} \int_0^{\infty} \frac{\varepsilon_p^{1/2} d\varepsilon_p}{e^{\varepsilon_p+\eta+\varepsilon_i+1}} = N_v F_{1/2}(-\eta - \varepsilon_i) \quad (4.48)$$

Here is the width of the zone where $E_c - E_0 = E_g$ - prohibited. The concentration of hives in this case is as follows:

$$N_v = 2 \left(\frac{2\pi m_p^* kT}{h^2} \right)^{3/2} \quad (4.49)$$

- density of effective cases in the valent zone;
- Fermi-Dirac integral for the valent zone.

If we hiosblaze that . This $m_r^* = m_p^* = m_0 T = 300$ K da $N_c = N_v = 2,5 \cdot 10^{19} \text{ sm}^{-3}$ means that there are 1/3 of 1 atom in 2000.

$$F_{1/2}(-\eta - \varepsilon_i) = \frac{2}{\sqrt{\pi}} \int_0^\infty \frac{\varepsilon_p^{1/2} d\varepsilon_p}{e^{\varepsilon_p + \eta + \varepsilon_i + 1}} \quad (4.50)$$

If $\frac{E_c - E_d}{kT} = \varepsilon_d$, if we introduce a definition, the concentration of electrons at the donor state based on (4.36) is:

$$n_d = f_0(E_d)N_d = \frac{N_d}{g^{-1}e^{-\varepsilon_d - \eta + 1}} \quad (4.51)$$

Concentration of positive ions in donor input:

$$p_d = f_{0p}(E_d)N_d = \frac{N_d}{ge^{\varepsilon_d + \eta + 1}} \quad (4.52)$$

We find the concentration of spikes and electrons in the inlet of the accorator. based on (4.38) and (4.39)

$$p_a = \frac{N_a}{g^{-1}e_a^{\varepsilon_a + \eta} + 1}; \quad (4.53)$$

$$n_a = \frac{N_a}{ge^{-\varepsilon_a - \eta} + 1}. \quad (4.54)$$

4.5. Semiconductor with input

To determine the equilibrium concentrations of electrons and swells in zones, it is necessary to calculate the Fermi-Dirac (4.44) and (4.50) integrals, which can be expressed as follows depending on the Fermi level cited

$-\infty < \eta < -1$ “When

$$F_{\frac{1}{2}}(\eta) = e^\eta \quad (4.55)$$

$-\infty < -\eta - \varepsilon_i < -1$ when if

$$F_{1/2}(-\eta - \varepsilon_i) = e^{-\eta - \varepsilon_i} \quad (4.56)$$

Such an approximation would be subject to Boltzman's statistics, while semiconductor would be undecomposed. For such a classical approximation, the concentration of balanced electrons and spurs

$$n_0 = N_c e^\eta; \quad (4.57)$$

$$p_0 = N_v e^{-\eta - \varepsilon_i}. \quad (4.58)$$

A classic approximation is shown in Fig. 4.5.

$5 < \eta < \infty$ When

$$F_{1/2}(\eta) = \frac{4}{3\sqrt{\pi}} \eta^{3/2} \quad (4.59)$$

$5 < -\eta - \varepsilon_{to} < \infty$ When

$$F_{1/2}(-\eta - \varepsilon_i) = \frac{4}{3\sqrt{\pi}}(-\eta - \varepsilon_i)^{3/2} \quad (4.60)$$

These approximations are appropriate for strongly articulated semiconductors. Fig. 4.5 shows for strongly coupled donor semiconductors

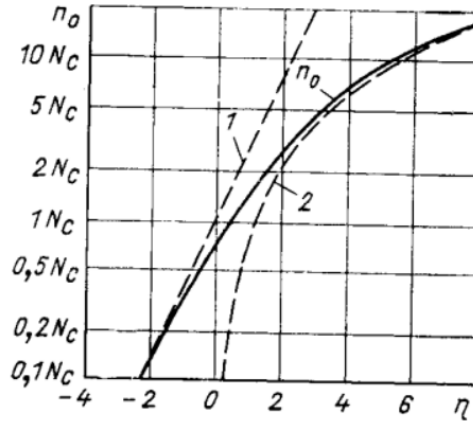


Fig. 4.5. The dependence of the concentration of conduction electrons on the quoted Fermi level

3. For the permeability zone when $\eta - 1 < < 5$

$$F_{1/2}(\eta) \approx \frac{1}{0,17 + e^\eta} \quad (4.61)$$

It can be written.

And for the valent zone when $\eta - 1 < -\varepsilon_i < 5$

$$F_{1/2}(-\eta - \varepsilon_i) = \frac{1}{0,27 + e^{-\eta - \varepsilon_i}} \quad (4.62)$$

Under such conditions, the semiconductor properties will be the conditions for the transition from the non-deformed state to the departed state. Major errors are made when calculating Fermi Dirac integrals by the formulas (4.61) and (4.62), so it is recommended to use special tables.

The n_0 curve in Fig. 4.5 (4.42) represents the concentration of balanced conduction electrons calculated by expression.

We'll look at borderline cases.

A semiconductor with an unmatched input. Based on Fig 4.5, the donor semiconductor can be considered non-decayed if

$$\eta = \frac{F - E_c}{kT} < -1$$

or

$$F < (E_c - kT). \quad (4.63)$$

If. This means that the unbroken donor is located at least kT below the conductivity zone in the Fermi level of the semiconductor in the prohibited zone. In this case, for electrons located in the conduction zone and having energy, the exponential gap in the Fermi-Dirac distribution would be much larger than once. So that $E \geq E_c$

$$f_0 = \frac{1}{e^{(E-F)/kT} + 1} \approx e^{-(E-F)/kT} \quad (4.64)$$

(4.64) states that charge carriers in unrelational donor semiconductors are subject to the Boltzman distribution. Then the Fermi-Dirac integral

$$F_{1/2}(\eta) = \frac{2}{\sqrt{\pi}} \int_0^\infty \frac{\varepsilon^{1/2} d\varepsilon}{e^{\varepsilon \dots \eta} + 1} = \frac{2}{\sqrt{\pi}} e^\eta \int_0^\infty e^{-\varepsilon} \varepsilon^{1/2} d\varepsilon = e^{-(E_c-F)/kT} \quad (4.65)$$

is written in the form. For an unfused semiconductor, the equivalent concentration of electrons and swells is respectively

$$n_0 = N_c F_{1/2}(\eta) = N_c e^\eta = N_c e^{-(E_c-F)/kT} \quad (4.66)$$

$$p_0 = N_v F_{1/2}(-\eta - \varepsilon_i) = N_v e^{-\eta - \varepsilon_i} = N_v e^{(E_v-F)/kT}. \quad (4.67)$$

The analysis of these relations shows that for the distribution of undischarged charge carriers, the Fermi integral coincides with the expression of the probability of filling the energy level E_c . It follows from (4.66) and (4.67), respectively, that in an unsplit input semiconductor, the concentration of electrons in the conductivity band will be the same as the concentration in the E_c state. The concentration of free spurs will be the same as putting the E_v level instead of the valent zone.

In a private semiconductor, free charge carriers are caused by the breakage of covalent bonds, so the concentration of free electrons and covalents is equal: $n_0 = p_0 = n_i$. n_i called the concentration of private charge carriers. If we define the Fermi level given in a private semiconductor η , based on (4.66) and (4.67)

$$n_i = N_c e^{\eta_i} = N_v e^{-\eta_i - \varepsilon_i}. \quad (4.68)$$

origin.

We now express through η the balanced concentration of electrons and swells for an undecomposed input semiconductor

$$\left. \begin{aligned} n_0 &= N_c e^\eta = n_i e^{\eta - \eta_i} \\ p_0 &= N_v e^{-\eta - \varepsilon_i} = n_i e^{\eta_i - \eta} \end{aligned} \right\}. \quad (4.69)$$

From here

$$n_0 p_0 = n_i^2 \quad (4.70)$$

origin.

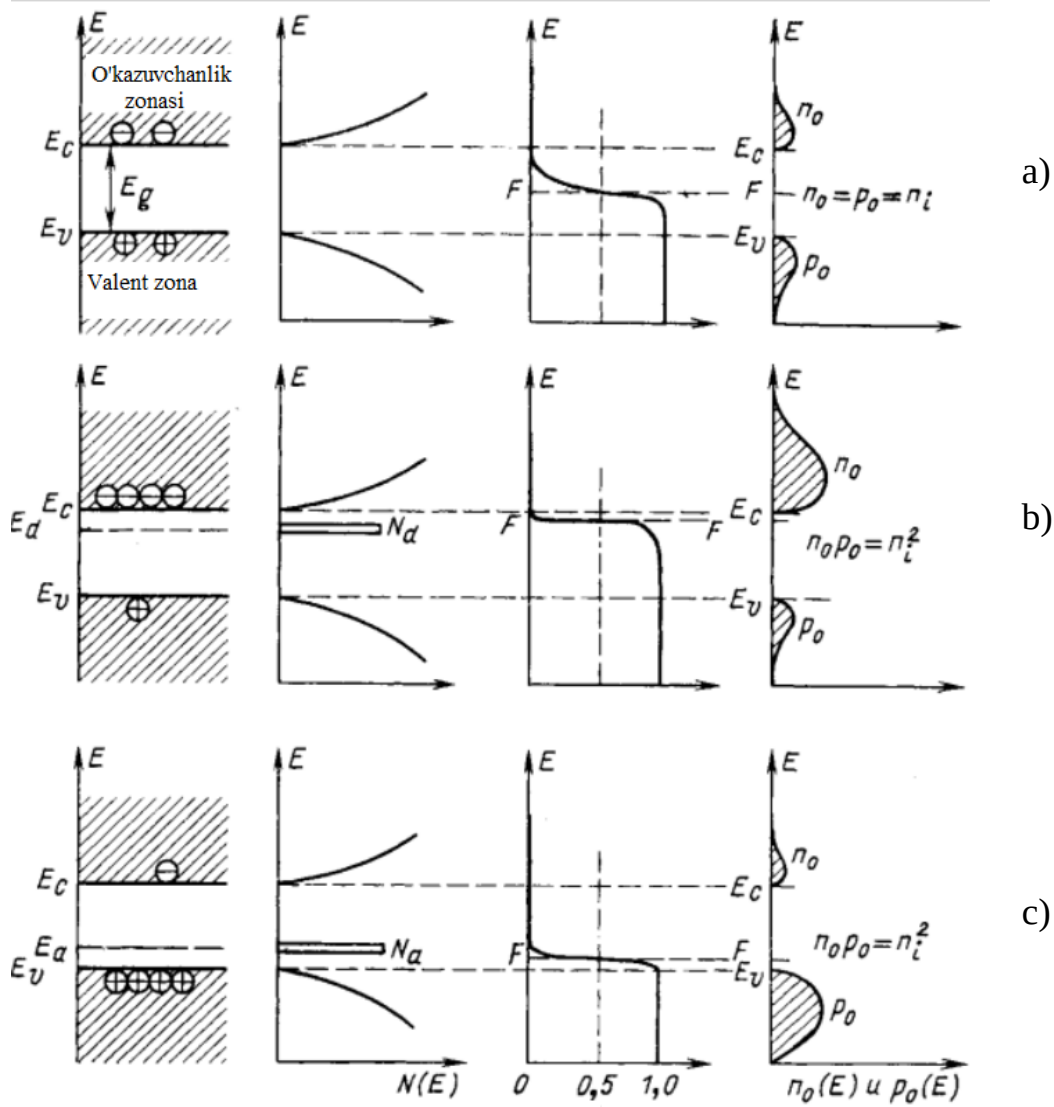


Fig. 4.6. Zone diagram in private, donor, and receiver semiconductors under the thermodynamic equilibrium state, state density, Fermi-Dirac distribution, and charge carrier concentration

Thus, in an unseparated input semiconductor, the multiplication of the concentrations of free electrons and spikes in thermodynamic equilibrium is a constant value equal to the square of the concentration of charge carriers in a particular semiconductor at a given temperature.

Figure 4.6 shows the zone diagram in private, electronic, and semiconductors in thermodynamic equilibrium, the density of the $N(E)$ states, the Fermi-Dirac distribution $f_0(E)$ and the concentration of charge carriers.

If heat excitation results in charge carriers, then the probability of such a process is proportional to $e^{-w/kT}$. The magnitude w is called the activation energy. It means that the electron concentration n_0 can also be expressed in terms of the activation energy

$$n_0 = Ce^{-w/kT}, \quad (4.70)$$

where C is constant.

(4.66) The attitude can be written as follows

$$n_0 T^{-3/2} = 2 \left(\frac{2\pi m_n^* k}{h^2} \right)^{3/2} e^{-(E_c - F)/kT} = C e^{-(E_c - F)/kT}. \quad (4.71)$$

From comparing this expression with (4.71)

$$w = -k \frac{d}{d\left(\frac{1}{T}\right)} \ln(n_0 T^{-3/2}) = E_c - F + T \frac{dF}{dT} \quad (4.72)$$

originates, that is, the activation energy depends on the Fermi level.

If there are free electrons in the conductivity zone at $T = 0$ K, then the semiconductor with the Fermi level given undergoes strong spin. In this case, the expression (4.42) is equal to the material and the equilibrium concentration of electrons in the strongly emitted donor semiconductor is equal to: $\eta = \frac{F - F_c}{kT} \rightarrow \infty$

$$n_0 = N_c \frac{2}{\sqrt{\pi}} \int_0^{\eta_0} \varepsilon^{1/2} d\varepsilon = \frac{4}{3\sqrt{\pi}} N_c \eta_0^{3/2} = \frac{8\pi}{3} N_c \left(\frac{2m_n^*}{h^2} \right)^{3/2} (F_0 - E_c)^{3/2} \quad (4.73)$$

It is determined by the concentration of balanced holes (4.67).

Fermi level at absolute zero based on (4.73)

$$F_0 = E_c + \frac{h^2}{2m_n^*} \left(\frac{3n_0}{8\pi} \right)^{2/3} \quad (4.74)$$

will take on a look. If and if, i.e., if the Fermi level is located below the valent band, the capacitor semiconductor will be strongly dissipated. The balanced concentration of free swells in a strongly rotated ejector semiconductor is based on (4.48) and (4.60) $(-\eta - \varepsilon_i) > 5F < (E_v - 5kT)5kT$

$$p_0 = N_v \frac{4}{3\sqrt{\pi}} (-\eta_0 - \varepsilon_i)^{3/2} = \frac{8\pi}{3} \left(\frac{2m_p^*}{h^2} \right)^{3/2} (E_v - F_0)^{3/2}, \quad (4.75)$$

It is determined from the concentration of balanced electrons (4.66).

From (4.73) and (4.75) it follows that the concentration of most charge carriers in a semiconductor with a strong oscillation does not depend on temperature. The Fermi level F_0 state is determined by the concentration of equilibrium charge carriers and depends on their effective mass. Moreover, the smaller the effective mass of the charge carriers, the smaller their concentration at which the spinning will occur.

Taking into account (4.74), the condition of complete relapse can be written as follows

$$\eta = \frac{F - E_c}{kT} = \frac{h^2}{2m_n^* kT} \left(\frac{3n_0}{8\pi} \right)^{2/3} \gg 1. \quad (4.76)$$

It is replaced here η_0 by η , which does not cause a significant error for a strong sobering state, as shown below. From the analysis (4.76), it follows that the earlier the spin occurs, the higher the concentration of free charge carriers, the lower the temperature (at a known concentration), and the smaller the effective mass of charge carriers.

Since T_n is the minimum concentration of charge carriers corresponding to the beginning of spin at a given temperature, we can write the following based on $F_0 = E_c + kT$ (4.73):

$$\begin{aligned} n_B &= \frac{4N_d}{3} \sqrt{\pi} = 3,78 \cdot 10^{15} \left(\frac{m_n^*}{m_0} \right)^{\frac{3}{2}} T_B^{\frac{3}{2}} = \\ &= 2 \cdot 10^{19} (m_n^*/m_0)^{3/2} (T_B/300)^{3/2} \end{aligned} \quad (4.77)$$

or

$$T_B = 4,1 \cdot 10^{-11} (m_0/m_n^*) n_B^{2/3}. \quad (4.78)$$

These conclusions are supported by experimental data. Fig. 4.7 shows the temperature-dependent curves of the concentration of silicon for silicon with different boron contents.

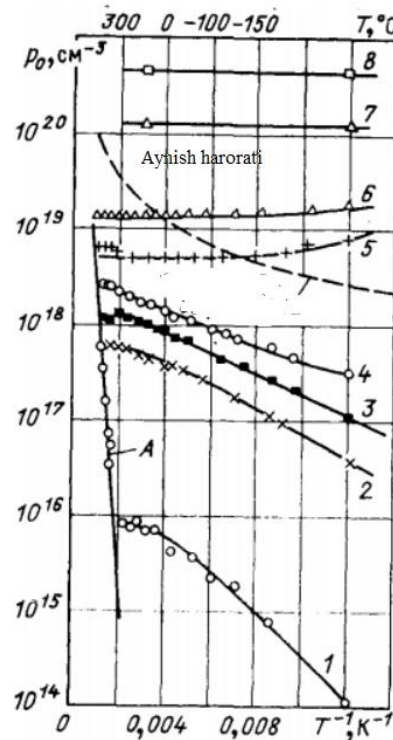


Fig. 4.7. Temperature dependence of concentration of hives for silicon alloyed with boron: A – specific concentration. Boron concentration: not legged; .1 – $2 - 6.7 \cdot 10^{17}$; 3 – $1.3 \cdot 10^{18}$; 4 – $2.7 \cdot 10^{18}$; 5 – $6.7 \cdot 10^{18}$; 6 – $1.3 \cdot 10^{19}$; 7 – $1.3 \cdot 10^{20}$; 8 – $1.3 \cdot 10^{21} \text{ sm}^{-3}$

The analysis of these dependencies shows that with an increase in the degree of oxidation, the concentration of the pit ceases to be dependent on temperature, that is, the semiconductor becomes strongly stained. The spin zone passes to a higher concentration of charge carriers as the temperature increases. If decomposition at liquid nitrogen temperature occurs at it, at room temperature it happens at U. Indium antimonide ($m^*=0.013m_0$), which has a very low effective electron mass at $T=300 \text{ K}$ $p_0 = 2.5 \cdot 10^{18} \text{ sm}^{-3}$ $p_0 = 2.0 \cdot 10^{19} \text{ sm}^{-3}$, turns out to be already decomposed at $n_0 = 2 \cdot 10^{16} \text{ sm}^{-3}$

For an isolated donor semiconductor, let's figure out the temperature-dependent Fermi level. If $\eta \gg 1$, then and the Fermi integral can be spread over the degree series and hence the equilibrium electron concentration $e^{\varepsilon-\eta} \ll 1$

$$n_0 = N_c \frac{4}{3\sqrt{\pi}} \eta^{3/2} \left(1 + \frac{\pi^2}{8\eta^2} - \frac{7\pi^4}{640\eta^4} + \frac{31\pi^6}{3072\eta^6} - \dots \right). \quad (4.79)$$

can be expressed by the formula.

Starting with the third, we throw out the hads and replace them with ni. In this, the error is somewhat reduced. For the result $\eta \eta_0 n_0$

$$n_0 \approx N_c \frac{4}{3\sqrt{\pi}} \eta^{3/2} \left(1 + \frac{\pi^2}{8\eta_0^2} \right). \quad (4.80)$$

We will have a review of it. (4.73) and (4.80)

$$\frac{3\sqrt{\pi} n_0}{4 N_c} = \eta^{3/2} \left(1 + \frac{\pi^2}{8\eta_0^2} \right) = \eta_0^{3/2}, \quad (4.81)$$

originates, from here

$$\eta = \eta_0 \left(1 + \frac{\pi^2}{8\eta_0^2} \right)^{-2/3}. \quad (4.81a)$$

$\eta_0 > 5$ when the second row in parentheses is small. (4.81a) Spread out in a row and limited to two hads,

$$\eta = \eta_0 \left(1 - \frac{\pi^2}{12\eta_0^2} \right) \quad (4.82)$$

or

$$F = F_0 \left[1 - \frac{\pi^2}{12} \left(\frac{kT}{F_0 - E_c} \right)^2 \right]. \quad (4.83)$$

we will have an expression. It follows from (4.83) that a semiconductor that has turned away at low enough temperatures F_0 will be very close to the Fermi level.

4.6. Private semi-digger

If the semiconductor penetration effect is small, then it is private. In it, free charge carriers occur only due to the breaking of the valent bonds (Fig. 4.8), so

$$n_0 = p_0 \quad (4.84)$$

The ratio (4.84) represents the electroneutrality condition of a special semiconductor, i.e., the resulting charge of all charged particles in the crystal must be zero. This requirement determines the state of the Fermi level. Uni

$$N_c F_{1/2}(\eta) = N_v F_{1/2}(-\eta - \varepsilon_i). \quad (4.85)$$

The equation can be deciphered.

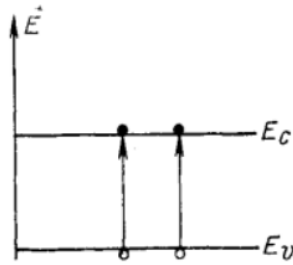


Fig. 4.8. Thermal generation of charge carriers in a private semiconductor

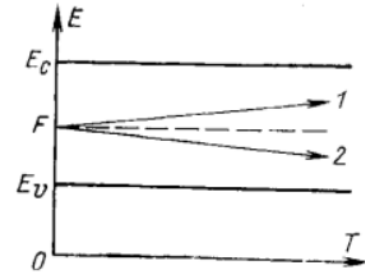


Fig. 4.9. Temperature dependence of Fermi temperature in a private semiconductor

A non-fused private semiconductor. For it, the zone diagram, the density of the cases, the Fermi-Dirac distribution, and the concentration of the charge carrier in thermodynamic equilibrium are shown in Fig. 4.6, a. For an uncoupled semiconductor, the Fermi-Dirac integral can be approximated by the exponential function (4.65). Bunda (4.85)

$$N_c e^\eta = N_t e^{-\eta - \varepsilon_i}, \quad (4.86)$$

takes the form, from where we determine the Fermi level

$$F = E_c - \frac{1}{2}E_g + \frac{kT}{2} \ln \frac{N_v}{N_c} = \frac{E_c + E_v}{2} + \frac{kT}{2} \ln \frac{N_v}{N_c}, \quad (4.87)$$

Here – the width of the forbidden zone. $E_g = E_c - E_v$

Agar (4.43) and N_c ni, (4.49) and N_v ni (4.87) ga qoysak,

$$F = \frac{E_c + E_v}{2} + \frac{3}{4} kT \ln \frac{m_p^*}{m_n^*}. \quad (4.88)$$

origin.

It follows from (4.88) that at absolute zero temperature, the Fermi level is located between the bottom of the permeability zone and the ceiling of the valent zone, i.e. . When it happens, its condition does not depend on the temperature and is placed in the middle of the prohibited zone. when the Fermi level shifts linearly to the bottom of the temperament zone as temperature increases, and when it does, it shifts linearly to the ceiling of the Fermi zone (Figure 4.9). $F = \frac{1}{2}(E_c + E_v)m_p^* = m_n^*m_p^* > m_n^*m_p^* < m_n^*$

We determine the concentration of private chargers

$$n_i = (n_0 p_0)^{1/2} = 4.9 \cdot 10^{15} \left(\frac{m_n^* m_p^*}{m_0^2} \right)^{3/4} T^{3/2} e^{-E_g/2kT} \quad (4.89)$$

Thus, the concentration of charge carriers in an undecomposed private semiconductor depends on the temperature T , the bandgap E_g width, the values of the m_n^* and m_p^* effective masses of the electrons, and does not depend on the state of the Fermi level. This is when the concentration is temperature-dependent $E_g \gg kT$ (4.89) is determined by the exponential had in it. This is confirmed by the data in Figure 4.10, which shows that the specific concentration of charge carriers in germanium, silicon, and gallium arsenide is dependent on the inverse temperature.

The activation energy of a private semiconductor follows from (3.72) and (3.87),

$$w = \frac{1}{2} E_g, \quad (4.90)$$

that is, half of the prohibited zone.

We evaluate the specific concentration of charge carriers in germanium, silicon and indium antimonide. For silicon. For German and . For indium antimonide, $m_{dn}^* = 1,08m_0, m_{dp}^* = 0,56m_0, T = 300 \text{ K}$ da $E_g = 1,11 \text{ eV}$, $n_i = 1,9 \cdot 10^{10} \text{ sm}^{-3}$. $1,11n_i = 1,9 \cdot 10^{10} \text{ sm}^{-3} m_{dn}^* = 0,56m_0, m_{dp}^* = 0,35m_0, E_g =$

$0,67 > Bn_i = 2 \cdot 10^{13} \text{ sm}^{-3} m_n^* = 0,013 m_0, m_{dp}^* = 0,6 m_0, E_g = 0,18 \text{ eV}, n_i = 2 \cdot 10^{16} \text{ sm}^{-3}.$

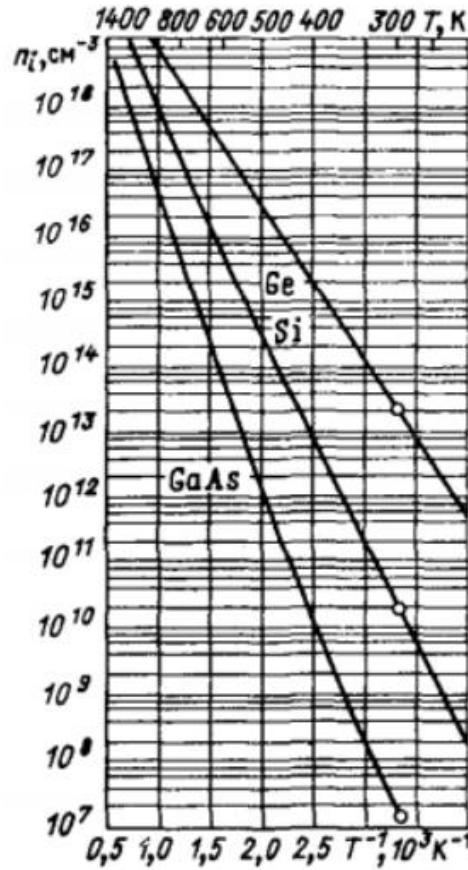


Fig. 4.10. Temperature dependence of concentration of private charge carriers in germanium, silicon and gallium arsenide

(4.89) is linearly linked to the function $\ln(n_i T^{-3/2}) 1/T$

$$\ln(n_i T^{-3/2}) = \text{const} - \frac{E_g}{2k} \frac{1}{T}. \quad (4.91)$$

The oblique angle of this line is $\text{tg } \theta = E_g/2k$. However, the data from Fig. 4.10 do not indicate that the E_g value does not correspond to the actual latitude of the German no-go zone. The reason for this discrepancy is that the forbidden zone itself changes with temperature. Increasing the amplitude of the thermal oscillations of the lattice atoms leads to its decrease. In addition, as the temperature increases, the interatomic distances change, which also affects the width of the forbidden zone. As a result, the $E_g(T)$ dependency can be of a complex nature. As an example, Fig. 4.11 shows the temperature-dependent changes in the width of the forbidden zone for germanium, silicon, and gallium arsenide. For these semiconductors, the bandgap values at 0 K are 0.89, 1.16, and 1.52 eV, respectively. They show that, as can be

seen from Fig. 4.11, the width of the exclusion zone in the temperature range of 175–350 K varies linearly with temperature. In this case, the change in the size of the zone is equal to and equal to the coefficient of temperature and for germanium, silicon, indium antimonide and gallyum arsenide, respectively $\alpha = dE_g/dT < 03,9 \cdot 10^{-4}; 2,4 \cdot 10^{-4}; 2,9 \cdot 10^{-4} 4,3 \cdot 10^{-4} \text{ V/K}$. For >0 in PbS α , it increases from 0.34 eV at 0 K to 0.41 eV at 300 K. In such cases, the dependency is subject to a linear law.

$$E_g = E_{g0} - \alpha T, \quad (4.92)$$

Here is the width of the prohibitory zone extrapolated at 0 K (Fig. 4.11, eV for GaAs). $E_{g0} E_{g0} = 1,58$

We put (4.92) into Equation (4.89)

$$n_i T^{-3/2} = 2 \left(\frac{2\pi^2 m_n^* m_p^* k^2}{h^2} \right)^{3/4} e^{-\alpha/2k} e^{-E_{g0}/2kT}, \quad (4.93)$$

from here

$$\ln(n_i T^{-3/2}) = \text{const} - \frac{\alpha}{2\hbar} - \frac{E_{g0}}{2k} \frac{1}{T}. \quad (4.94)$$

We determine it

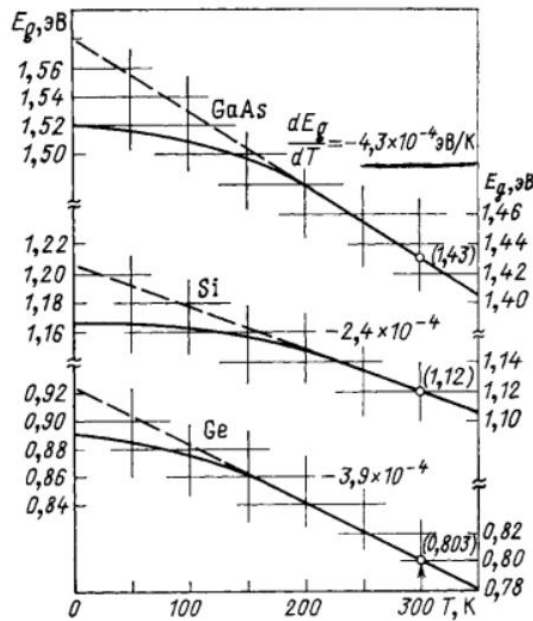


Fig. 4.11. Temperature of the width of the forbidden zone in germanium, silicon and gallium arsenide

Thus, if the forbidden zone width of a semiconductor is linearly dependent on temperature, then the graph of $\ln \left(n_i T^{-\frac{3}{2}} \right) 1/T$ is also a straight line, whose slope is characterized by its value of E_{go} , the width of the banned zone extrapolated at 0 K. The actual value of the bandwidth of a semiconductor is determined by the formula (4.29) at a given temperature.

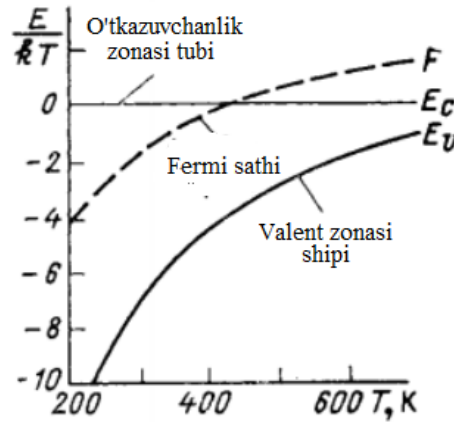


Fig. 4.12. Temperature dependence of Fermi temperature for indium antimonide

A semiconductor. According to (4.88) for a private semiconductor, the rate of change of the Fermi level with temperature is proportional to the ratio of the efficacy masses of the swells and electrons. As a result, with an increase in temperature, the Fermi level moves away from the zone of heavy charge carriers, approaching the zone of light charge carriers. If the distance from the Fermi level to this zone corresponds to the value of kT , then the rotation occurs, and the corresponding Fermi-Dirac integral is no longer replaced by an exponential. Moreover, the more the effective masses of electrons and canines differ, the faster the spin occurs. If, for example, a distortion occurs in the permeability zone, it is absent in the tensile zone, and the expression of the concentration of the individual carriers takes the following form:

$$n_i = N_c F_{1/2}(\eta) = N_v e^{-\eta - \varepsilon_i} \quad (4.95)$$

From here, we get the following formula to determine the state of the Fermi level

$$e^{\eta} F_{1/2}(\eta) = (m_p^*/m_n^*)^{3/2} e^{-\varepsilon_i} \quad (4.96)$$

A classic example of semiconductors with strongly varying effective masses is indium angimonide, for which . Fig. 4.12 shows the change in the state of the Fermi level in InSb depending on temperature. In this case, the dependence is

considered, which, according to experimental data, has the following form: $m_p^* \approx 10m_t^*E_g(T)$

$$E_g = (0,26 - 2,7 \cdot 10^{-4}T) \text{ eV.}$$

As can be seen from Fig. 4.12, at $T=200 \text{ K}$, the Fermi level lies only $4kT$ below E_c , and at about 450 K it passes through the E_c level and enters the permeability zone, which causes a strong spin of the electron gas. This must be taken into account, and parity must be used in calculating (4.95).

4.7. Inlet concentration and temperature dependence of the Fermi level for an unfused semiconductor

A semiconductor can contain a donor and an acceptor input. The concentrations of these inputs are respectively and be them. N_d and N_a

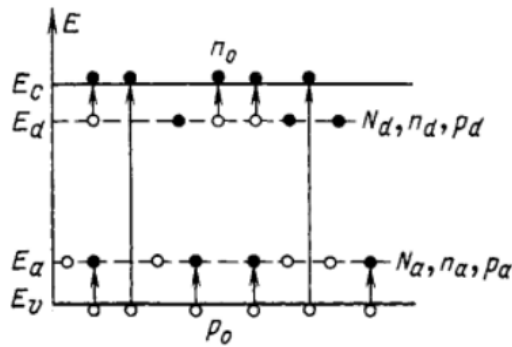


Fig. 4.13. Donor and Acceptor Input Semiconductor

The energy diagram of such a semiconductor is shown in Fig. 4.13. Spikes in the conduction zone and cavities in the stress band are formed by thermal excitation of electrons from the tense band and incoming atoms. These transitions are shown with arrows in Fig. 4.13. To determine the state of the Fermi state in a semiconductor in thermodynamic equilibrium, the electron neutrality condition can be used, according to which the total charge of all charged particles at any point in the crystal must be zero. For a semiconductor with two types of inputs, the electron neutrality equation is written as follows

$$n_0 + n_a = p_0 + p_d \quad (4.97)$$

where $n_a = N_a - p_u$ is the number of electrons connected to the accessory of the acceptor; $-$ the number of ionized atoms of donor inputs. and using expressions for (4.97), the condition of electroneutrality can be written as follows $p_d = N_d - n_d n_a p_d$

$$(n_0 + n_d) - (p_0 + p_u) = N_d - N_a. \quad (4.98)$$

We use expressions (4.42), (4.48), (4.51), and (4.53)

$$N_c F_{1/2}(\eta) + \frac{N_d}{g^{-1}e^{-\varepsilon_d-\eta}+1} - N_v F_{1/2}(-\eta - \varepsilon_i) - \frac{N_a}{g^{-1}e^{\varepsilon_a+\eta}+1} = N_d - N_a. \quad (4.99)$$

For non-fused semiconductors (4.99), parity is much simpler

$$N_c e^\eta + \frac{N_d}{g^{-1}e^{-\varepsilon_d-\eta}+1} - N_v e^{-\eta-\varepsilon_i} - \frac{N_a}{g^{-1}e^{\varepsilon_a+\eta}+1} = N_d - N_a, \quad (4.100)$$

However, even in this case, determining the Fermi level faces great difficulties. Therefore, in the future, we will consider the solution of special problems of great practical importance.

1. Donor semiconductor ($N_a = 0$)

First of all, we are limited only to the temperature range in which ionization of the input centers occurs and the internal conductivity is absent, i.e. $p_0 = 0$. The condition of electroneutrality (4.98) reads as follows:

$$n_0 + n_d = N_d \quad (4.101)$$

or

$$n_0 = p_d$$

based on (4.52)

$$n_0 = \frac{N_d}{ge^{\varepsilon_d-\eta}+1} \quad (4.102)$$

It's a we get

Let's assume that there is no spinning. Without . From here $n_0 = N_c e^\eta$

$$e^\eta = n_0/N_c. \quad (4.103)$$

We clarify that By subtracting (4.103) by (4.102), we get the following equation for the concentration of electrons in the conduction band

$$n_0 = \frac{N_c N_d}{N_c + ge^{\varepsilon_d} n_0}, \quad (4.104)$$

We form the following equation

$$n_0^2 + g^{-1}e^{-\varepsilon_d}N_c n_0 - g^{-1}e^{-\varepsilon_d}N_c N_d = 0 \quad (4.105)$$

We define this by solving the square equation n_0

$$n_0 = \frac{g^{-1}e^{-\varepsilon}d_{N_c}}{2} \left(\sqrt{V1 + 4ge^{\varepsilon}N_c^{-1}N_d} - 1 \right) \quad (4.106)$$

We illustrate this relationship as follows

$$n_0 = \frac{2N_d}{V + 4ge^{\varepsilon}N_c^{-1}N_d + 1}. \quad (4.107)$$

The Fermi level can be determined from (4.103)

$$e^{\eta} = \frac{n_0}{N_c} = \frac{2N_d/N_c}{1 + 4ge^{\varepsilon}N_c^{-1}N_d + 1} \quad (4.108)$$

We logarithm this equation

$$\eta = -\ln \left[\frac{N_c}{2N_d} + \sqrt{\left(\frac{N_c}{2N_d} \right)^2 + ge^{\varepsilon}N_cN_d^{-1}} \right] \quad (4.109)$$

or

$$F = E_c - kT \ln \left[\frac{N_c}{2N_d} + \sqrt{\left(\frac{N_c}{2N_d} \right)^2 + ge^{\varepsilon}N_cN_d^{-1}} \right] \dots \quad (4.110)$$

At lower temperatures, when (4.110) the formula is simplified and the Fermi level $e^{(E_c-E_d)/kT} \gg N_c/2N_d$

$$F = E_c - \frac{E_c - E_d}{2} - \frac{kT}{2} \ln \frac{gN_c}{N_d} = \frac{E_c + E_d}{2} - \frac{kT}{2} \ln \frac{gN_c}{N_d}, \quad (4.111)$$

and the activation energy is based on (4.72)

$$w = \frac{1}{2}(E_c - E_d),$$

It turns out that the donor input is equal to half the ionization energy.

Based on the (4.111) relationship, in the donor semiconductor, at an absolute zero temperature, the Fermi level is located between the bottom of the conduction zone and the donor input surface. In the region of sufficiently low temperatures (a few degrees on the Kelvin scale), when $gN_c < N_d$, the Fermi level first rises to some kind of maximum value, and then begins to decrease, and when $gN_c = N_d$, it becomes equal to again, as in the case of $T=0$ $F = 1/2 (E_c + E_d)$ A further rise in temperature is accompanied by an increase in N_c , and when there is $gN_c = N_d$ in the thermal degradation area, the Fermi level will continue to decline. Such a shift in

the Fermi level coincides with the indicate dependence of the concentration of electrons on the temperature

$$n_0 = N_c e^{\eta} = N_c e^{-\frac{(E_c - E_d)}{2kT} - \frac{1}{2} \ln \left(\frac{g N_c}{N_d} \right)} = (g^{-1} N_c N_d)^{\frac{1}{2}} e^{-\frac{(E_c - E_d)}{2kT}} \quad (4.112)$$

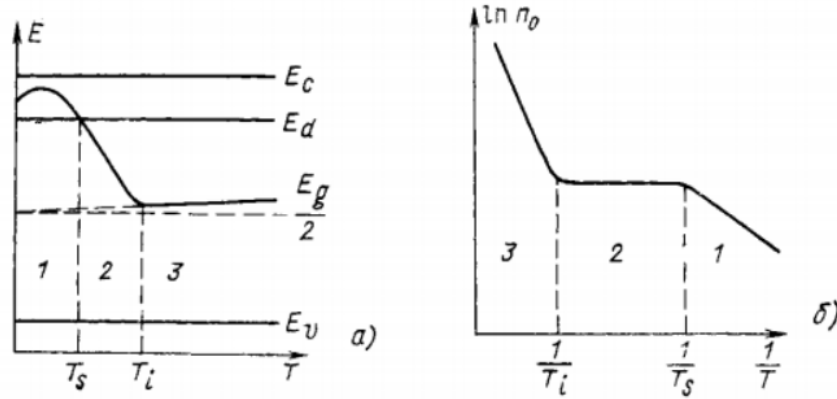


Fig. 4.14. Fermi level and temperature dependence of electron concentration for a donor semiconductor

The area of change in the Fermi level with the temperature described in the formula (4.111) is the weak ionization region (or freezing region) of the mixture. It is denoted by the number 1 in Fig. 4.14, which shows the variation in Fermi level and electron concentration, which is temperature-dependent for the donor semiconductor.

With a further increase in temperature, the concentration of electrons in the conductivity zone becomes comparable to the input concentration, and in this case the expressions (4.111) and (4.112) will not apply. However, now we can consider another extreme case, if the temperature is high enough and

$$e^{(E_c - E_d)/kT} \ll N_c / 2N_d$$

if inequality persists. In which (4.110) the expression is approxified as follows

$$F = E_c - kT \ln(N_c / N_d) \quad (4.113)$$

It means

$$n_0 = N_d \quad (4.114)$$

This means that almost all donor inputs are ionized and the concentration of electrons in the conduction zone is not dependent on temperature. This area of temperature in which the full ionization of the input occurs is called the reduction of the inlet (or the region of full ionization of the inlet) and is denoted by the number 2 shown in Fig. 4.14. When the complete ionization condition of the donor mixture is $n_0 = N_d$, E_T corresponds to the Fermi level below several levels of the mixture. This

means that as temperatures increase, the Fermi level drops, passes through the E_d level and goes down. The temperature at which $F=E_d$ is called the saturation temperature T_s , which

$$F = E_c - kT_s \ln(N_c/N_d) = E_d, \quad (4.115)$$

can be determined from the condition. From here

$$T_s = \frac{E_c - E_d}{k \ln(N_c/N_d)} \quad (4.116)$$

As can be seen from (4.116), the lower the decay temperature, the higher the effective electron mass, which determines the value of ionization energy (E_c-E_d) and the donor input concentration N_{dj} and N_c . Reduction of pollution at small values (E_c-E_d) occurs at very low temperatures. For example, in an $N_d = 10^{16} \text{ sm}^{-3}$ electron germanium with a sweep of, the ionization energy is 0.0096 eV, while saturation occurs at $T_s=32 \text{ K}$.

With a further increase in temperature, the concentration of the electron in the conductivity zone increases due to the electron transitions through the stress zone. In this, the state of the Fermi level and the electron concentration are determined by equations (4.88), (4.89) and (4.115). In Fig. 4.14, the 3 spheres correspond to the sphere of private conductivity. In this case, $F=F_i$ and the following expression from (4.88) and (4.113)

$$F_i = E_c - \frac{1}{2}E_g - \frac{kT_i}{2} \ln \frac{N_c}{N_0} = E_e - kT_i \ln \frac{N_d}{N_c}. \quad (4.117)$$

where

$$T_i = \frac{E_g}{k \ln \left(\frac{N_c N_0}{N_d^2} \right)} \quad (4.118)$$

Analysis of this expression shows that the lower the temperature T_i at which internal conductivity occurs in the donor semiconductor, the lower the bandwidth and inlet concentration, and the higher the value of the effective masses of the charge carriers.

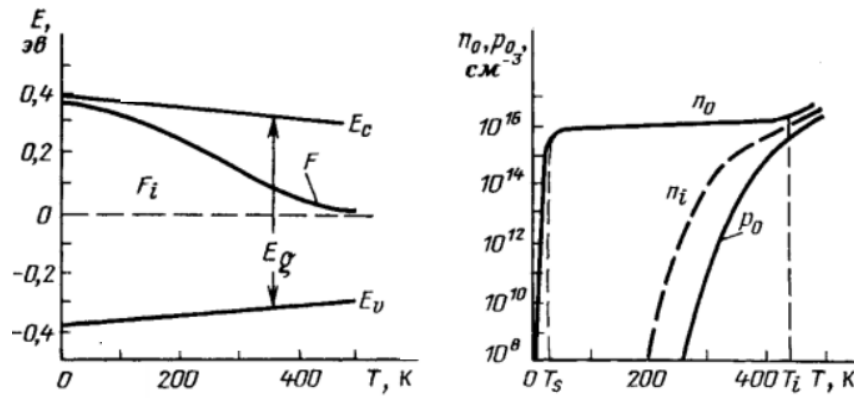


Fig. 4.15. Fermi level and temperature dependence of charge carrier concentration for germanium alloyed by surge.

Thus, using the approximations described, it is possible to observe the change in electron concentration and the state of the Fermi level in the electron semiconductor's forbidden zone over the entire temperature range.

As an example, the temperature dependencies of the Fermi level in Fig. 4.15 a and the concentration of the equilibrium electrons in Figure 4.15 are n_0 and the sum of p_0 for germanium added by surge. In addition, the dot line in these curves $N_d \approx 10^{16} \text{ cm}^{-3}$ indicates the rate of F_i and n_i in local German. When compiling the graphs, the temperature dependence of the German exclusion zone was taken into account.

From this figure it follows that at absolute zero temperature, the Fermi level in germ is located between the bottom of the E_c permeability zone and the donor input level E_d . As the temperature rises, it drops and the inlet level approaches E_d . At saturation temperatures, electrons are present in an equal amount to the donor input:

$$n_d = \frac{N_d}{\frac{1}{2} e^{(E_d - F)/kT} + 1} = \frac{2}{3} N_d$$

and in the conduction zone, respectively, $1/3 * N_d$ electrons. As temperatures increase further, the Fermi level will continue to decline, and the area of decline will begin; all inputs are ionized and the concentration of conduction electrons remains constant and is equal to $n_0 = N_d$. It is in this region of temperature that ionization of atoms of the main substance occurs, and non-main charge carriers appear - canines. Their concentration is as follows

$$p_0 = \frac{n_i^2}{n_0} = \frac{N_c N_v}{N_d} e^{-E_g/kT}. \quad (4.119)$$

In response, it increases sharply with an increase in temperature.

When the Fermi level reaches the middle of the no-go zone, then $n_0=p_0=n_i$ and the semiconductor transitions from the inlet to the interior. With a further increase in temperature, the Fermi level approaches the zone with a low effective density of cases.

2. Acceptor semiconductor ($N_d = 0$)

The determination of the Fermi level in an unbroken semiconductor that is alloyed with an acceptor input is the same as that of a donor semiconductor. The condition of electroneutrality is written as:

$$p_0 = n_a = N_a - p_a. \quad (4.120)$$

That's for the non-decoupled semiconductor, considering (4.67) and (4.54)

$$p_0 = \frac{N_a}{ge^{-\varepsilon} a^{-\eta} + 1} = \frac{N_a N_v}{N_v + ge^{\varepsilon_i - \varepsilon_a} \rho_0}. \quad (4.121)$$

After not so difficult algebraic modifications

$$p_0 = \frac{2N_a}{\sqrt{1 + 4ge^{\varepsilon_i - \varepsilon_a} N_a N_v^{-1}} + 1}. \quad (4.122)$$

We take it. Fermi level status

$$e^{\eta + \varepsilon_i} = \frac{N_v}{p_0} = \frac{N_v}{2N_a} \left[\sqrt{1 + 4ge^{\varepsilon_i - \varepsilon_a} N_a N_v^{-1}} + 1 \right]. \quad (4.123)$$

is determined by equality. By logarithming both sides of this equation,

$$F = E_c - E_g + kT \ln \left[\frac{N_v}{2N_a} + \sqrt{\left(\frac{N_v}{2N_a} \right)^2 + gN_v N_a^{-1} e^{(E_a - E_v)/kT}} \right] \quad (4.124)$$

We will have it. Expression at small temperatures, when $e^{(E_a - E_0)/kT} \geq \frac{N_v}{2N_a}$

$$F = E_c - E_g + \frac{1}{2}(E_a - E_v) + \frac{kT}{2} \ln \left(\frac{gN_0}{N_a} \right) = \frac{E_v + E_a}{2} + \frac{kT}{2} \ln \left(\frac{gN_0}{N_a} \right) \quad (4.125)$$

comes to the view.

Thus, in a hollow semiconductor, the Fermi level at absolute zero is positioned between the valent zone ceiling and the envelope intake surface. As the temperature increases, the Fermi level decreases when it occurs, then rises to the level. The concentration of free hives in this area = $\frac{1}{2}(E_0 + E_u)N_a > gN_v E_a$

$$p_0 = N_v e^{-\eta - \varepsilon_t} = (g^{-1} N_u N_v)^{1/2} e^{(E_0 - E_a)/2kT}. \quad (4.126)$$

It is determined by attitude. At high temperatures, when (4.124) $e^{(E_a - E_v)/kT} \ll \frac{N_0}{2N_a}$

$$F = E_c - E_g + kT \ln \frac{N_0}{N_a} \quad (4.127)$$

Concentration of free hives

$$p_0 = N_v e^{-\ln(N_v/N_a)} = N_a. \quad (4.128)$$

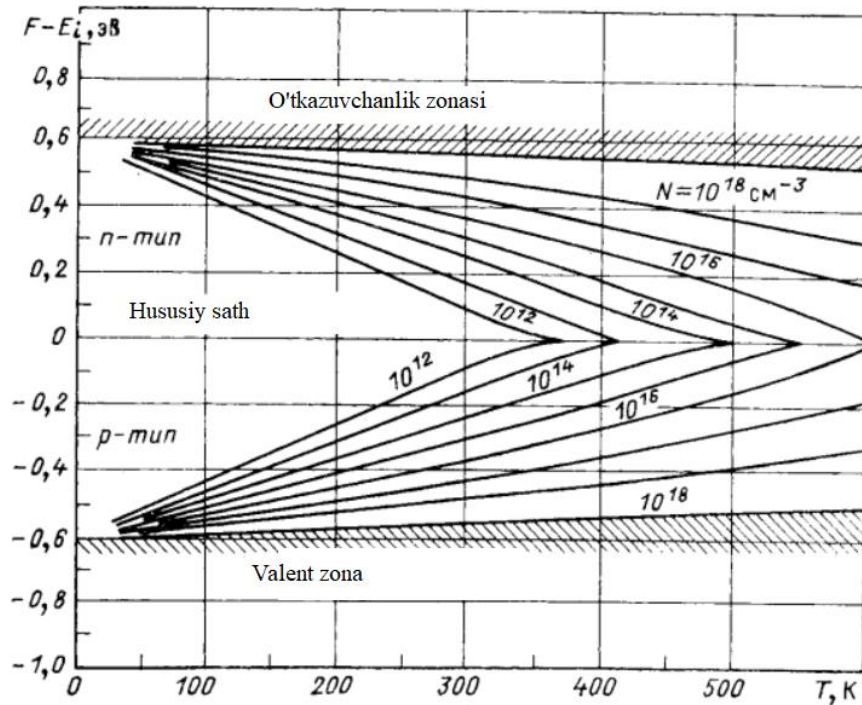


Fig. 4.16. The dependence of Fermi levels in silicon on temperature and concentration of inputs

In the acceptic semiconductor, as in the donor inlet, a poorning region is formed, characterized by the complete ionization of the atoms of the acceptor inlet at high temperatures. With a further increase in temperature, the Fermi level rises to the middle of the exclusion zone, and the semiconductor behaves as if it were private.

The Fermi level for silicon depending on the inlet concentration and temperature is shown in Fig. 4.16. Here is the temperature dependence of the no-go zone.

4.8. Temperature-dependent Fermi level for a partially compensated unseparated semiconductor with input

The considered cases of pure donor and pure ejector semiconductor are generally not applicable in practice. No matter how thoroughly the material has been cleaned, there will usually be more than one type of mixture left in it, and some entry centers may donor receive more than one electron.

In addition to the single-electron donor input, we will dwell in more detail on the state in which the semiconductor has a certain number of acceptor mixtures (see Fig. 4.13).

Suppose that in a semiconductor let the donor and receptor mixtures be equal quantities, i.e., $N_d = N_a$. At absolute zero temperature, electrons occupy the lowest energy states. This means that there are no electrons in the conduction band, and all cases are located in the stress band. Because there are free states in the acceptic input that are energetically lower than the donor input states, the electrons move from the donor state to the acceptor state, trying to fill the states with the lowest energy. As a result, positive and negative ions are formed in equal quantities - from which the electric neutrality condition is satisfied. Such a semiconductor is called compensated. In it, the mutual compensation of mixtures takes place, as a result of which they cannot become suppliers of electrons and hives; Thermal formation of charge carriers occurs as a result of the "zone zone". $N_d^+ = N_a^-$

In a compensated semiconductor, because of the presence of mixtures, the circuiticity of the lattice field is distorted much more strongly than in a true specific semiconductor. This affects the effects associated with the dispersion of charge carriers, in particular, their mobility.

We now assume that the Fermi level over the entire temperature range of $N_d > N_a$ and the entire temperature under consideration is significantly higher than E_a , i.e., higher than the energy level of the receiving input. In this case, all the acceptor states are filled with electrons transferred from the donor input layer. As a result, the total number of electrons on the donor layers and in the conduction zone is equal to $N_d - N_a$. There will be no empty hives in the tensile zone until special conductivity occurs, and the semiconductor will behave as a net donor. The same semiconductor is called partially compensated. The level of compensation is determined by the ratio of $N_d - N_a$. In such a semiconductor, the transfer of charge is carried out only by electrons, and the Hall constant has a negative direction.

If the N_a/N_d and Fermi levels are significantly lower than E_d , the donor input is fully ionized. All electrons have passed from the donor level to the recipient input layers, so that at $T = 0$, N_a/N_d hives appear in the receptor input atoms. As the temperature increases, electrons from the stress zone pass into the receiving input atoms, and the semiconductor behaves like a hollow semiconductor with a positive Hall coefficient. With the onset of internal conductivity, the gesture of the Hall

constant changes, since the mobility of electrons is greater than the mobility of the holes.

For the N_d/N_a state, we determine the temperature dependence of the concentration of free electrons in a semiconductor with a partially compensated mixture and the Fermi condition. For ionizing only the input, the electric neutrality condition is of the following form:

$$n_0 + n_d = N_d - N_a \quad (4.129)$$

$$n_0 + N_a = p_d. \quad (4.130)$$

Let's assume that the semiconductor is uneven. using (4.52) and (4.130) $n_0 = N_c e^\eta$

$$n_0 + N_a = \frac{N_d N_c}{N_c + g e^{\varepsilon_d} n_0} \quad (4.131)$$

We get the expression. Solve this equation,

$$n_0 = \frac{2(N_d - N_a)}{1 + \frac{g N_a}{N_c} e^{\varepsilon_d} + \left[\left(1 + \frac{g N_a}{N_c} e^{\varepsilon_d} \right)^2 + \frac{4g(N_d - N_a)e^{\varepsilon_d}}{N_c} \right]^{1/2}}. \quad (4.132)$$

we will get the result. From the relationship we determine the Fermi energy $e^\eta = n_0/N_c$

$$F = E_c + kT \cdot \left[\frac{4g(N_d - N_a)e^{(E_c - E_d)/kT}}{N_c} \right]^{1/2} \quad (4.133)$$

As can be seen, this result is consistent with the formulas (4.107) and (4.110) deduced for the undeformed semiconductor. $N_a = 0$

In the region of high temperatures, when the donor atoms are fully ionized

$$\frac{g N_a}{N_c} e^{\varepsilon_d} \ll 1 \text{ и } \frac{4g(N_d - N_a)e^{\varepsilon_d}}{N_c} \ll 1. \quad (4.134)$$

Inequality is appropriate. These inequalities

$$n_0 = N_d - N_a \quad (4.135)$$

$$F = E_c + kT \ln \frac{N_d - N_a}{N_c} \quad (4.136)$$

allows you to get relationships.

As expected, the concentration of the free electrons is constant and is determined by the amount of uncompensated donor access. The Fermi level

decreases with increasing temperatures, approaching the middle of the prohibited zone at temperatures corresponding to the onset of private permeability.

In the zone of low temperatures, the form of expression for n_0 and F depends on the ratio of the concentration of donors and receptors. If the degree of compensation is small, i.e., $N_a \ll N_d$, then it is possible to find a temperature range in which the ratio $N_a \ll n_0 \ll N_d$ which is equivalent to the condition is satisfied.

$$\frac{gN_a}{N_c} e^{\varepsilon_d} \ll 1 \ll \frac{4g(N_d - N_a)e^{\varepsilon_d}}{N_c}. \quad (4.137)$$

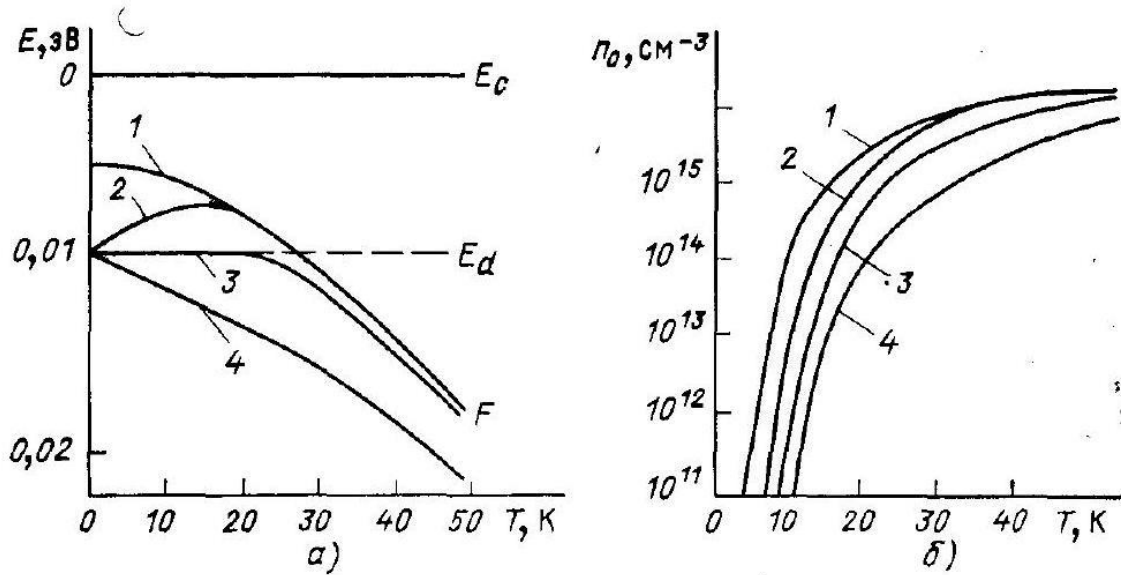


Fig. 4.17. The temperature dependence of the Fermi level and electron concentration at different compensation levels in Germanium. : $N_d - N_a = 10^{16} \text{ sm}^{-3}$. N_a 1 – 0; 2 – 10^{14} ; 3 – 10^{16} ; 4 – 10^{18} sm^{-3}

The concentration and Fermi levels based on (4.132) and (4.133) in this area are characterized by the following interpretation formulas

$$n_0 \approx (g^{-1}N_cN_d)^{1/2} e^{-(E_c - E_d)/2kT}; \quad (4.138)$$

$$F = \frac{E_c + E_d}{2} + \frac{kT}{2} \ln \frac{N_d}{gN_c}, \quad (4.139)$$

They are consistent with the expressions (4.112) and (4.111) for donor semiconductors. It turns out that the concentration of free electrons at the same time does not depend on the number of atoms in the receiver. $N_d \gg N_a n_0$

At $T = 0$, the Fermi level is located between E_c and E_d . As the temperature increases, it decreases (4.17, curve 1 in Fig. a). However, at lower temperatures, when $n_0 \ll N_a$, the effect of the receivers can no longer be ignored. The state of Fermi levels in this area is significantly different from what happens in a pure donor semiconductor.

In the case of strong compensation, inequality is appropriate, so from the expressions (4.132) and (4.133) $n_0 \lll N_a < N_d \frac{gN_a}{N_c} e^{\varepsilon_d} \gg 1 \frac{gN_a}{N_c} e^{\varepsilon_d} > \frac{4g(N_d-N_a)e^{\varepsilon_d}}{N_c}$

$$n_0 \approx N_c \frac{N_d - N_a}{gN_a} e^{-(E_c-E_d)/kT}; \quad (4.140)$$

$$F = E_d + kT \ln \frac{N_d - N_a}{gN_a} = E_d + kT \ln \frac{N_d - N_a}{2N_a}, \quad (4.141)$$

arises because $g = 2$.

In this case, the activation energy $w=E_c-E_d$ is equal to the ionization energy of the donor input, i.e., 2 times that of the pure donor semiconductor. Based on the relationship in Fig. 4.17 (4.141), the $F(T)$ relationship is plotted, and Fig. 4.17-b shows the $n_0(T)$ relationship for different levels of compensation. It can be seen that as the temperature increases, the absolute starts from zero, if any, the Fermi level increases first and then decreases below the donor entry level (curve 2). $\frac{N_d-N_a}{2N_a} > 1$ In the opposite case, as the temperature increases, the Fermi level decreases from the donor level (curve 4), and when it does, it remains unchanged until the condition $n \frac{N_d-N_a}{2N_a} < 1 \frac{N_d-N_a}{2N_a} = n_0 \lll N_a$ is met (curve, then decreases). From the comparison of the curves shown in Fig. 4.17,b, it follows that the temperature at which the $n_0(T)$ dependency is saturated increases with the increase in the compensation level of the donor input.

4.9. Semiconductors with access at extremely low temperatures

We highlight another property of semiconductors with inputs. At temperatures insufficient for ionization of the main substance in the donor semiconductor, the presence of electrons in the conduction zone is determined by thermal leaching of the donor input atoms. In this case, the number of conduction electrons is equal to the number of positive input ions formed, i.e.

$$n_0 = p_d = \frac{N_d}{g e^{(F-E_d)/kT} + 1} \quad (4.142)$$

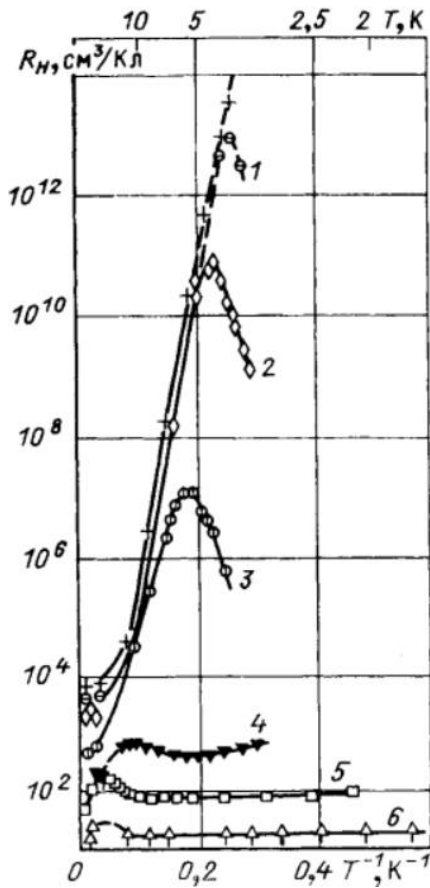


Fig. 4.18. Temperature dependence of the Hall constant for p-germanium

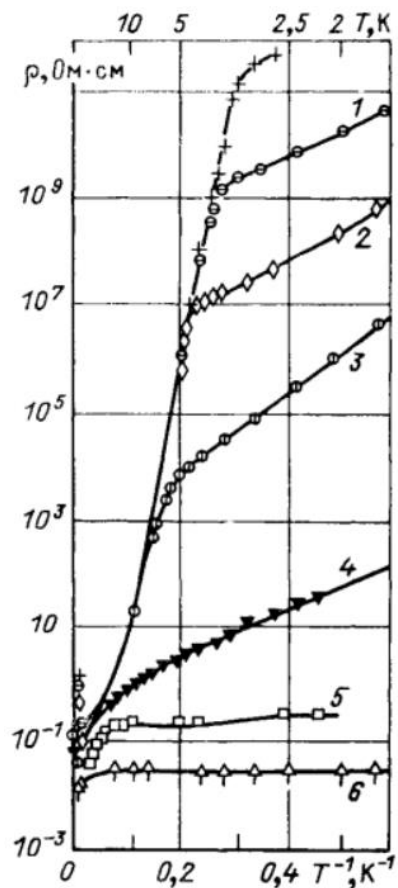


Fig. 4.19. Temperature dependence of specific resistance for P-germanium

In such a semiconductor, the input should increase exponentially with the specific resistance of the semiconductor and the Hall coefficient $\rho = (en_0\mu)^{-1}$ according to temperatures where the kT is less than the ionization energy of the donor or receiver (4.142). However, studies have shown that for semiconductors with input concentrations above 10^{17} cm^{-3} , the Hall coefficient exceeds maximum at lower temperatures. The state of this maxim depends on the input concentration, as can be seen from Fig. 4.18. The dependence of resistance of these samples on temperature is shown in Fig. 4.19. At temperatures where the maximum Hall coefficient occurs, resistance reaches saturation and increases more slowly with further temperature drops. For samples with input content $\delta 10^{19}$ higher than cm, resistance does not depend on temperature.

From these experimental facts, it follows that at low temperatures, a phenomenon called access band permeability occurs in a semiconductor with a high input content. In this case, the mechanism of input conductivity differs depending on the input concentration introduced.

The hydrogen-like model of input atoms does not account for their interactions with each other. However, at a sufficiently high concentration of input

atoms, given their interaction, it will be necessary because, according to (2.175), the radius of Boron for an electron bonded at a hydrogen-like center of entry is proportional to the relative dielectric conductivity and inversely proportional to the effective mass of the electron. Therefore, the effective orbital radius for the ingress electron in germanium is about 60 times greater than the Boron radius of an electron in a hydrogen atom. Therefore, even at the input concentration that is approx. 100, the wave functions of the electrons bonded to the neutral donor mixture are significantly consistent. 10^{16} sm^{-3}

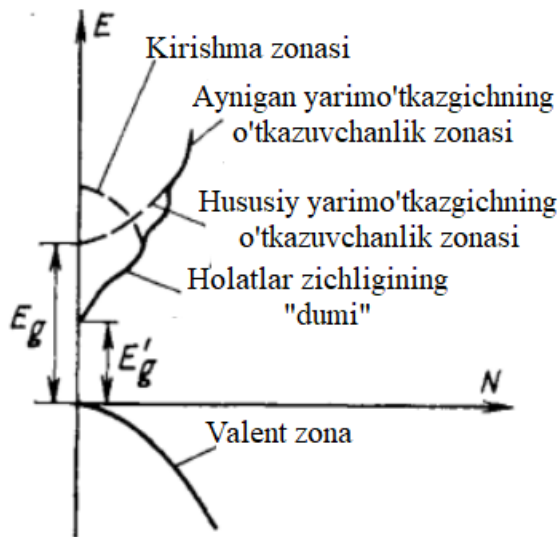


Fig. 4.20. Case density for fused semiconductor

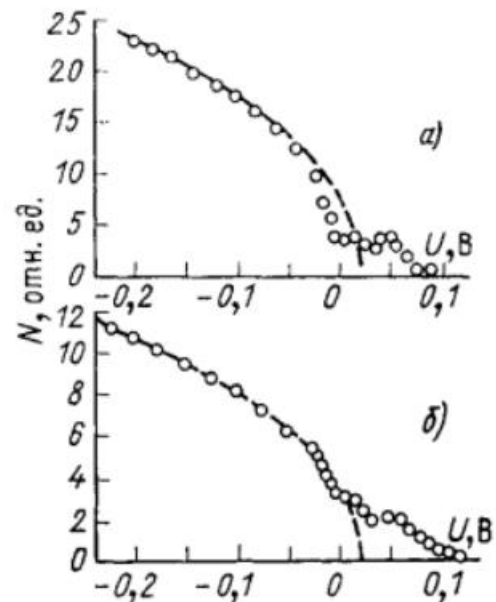


Fig. 4.21. Case density for P-type gallium arsenide

The width of this energy field depends on the degree to which the wave functions of the input atoms adhere to each other, and therefore depends on the donor input concentration. In this case, the conductivity occurs through the intake zone, which is caused by the electrons of the donor input, which behaves like an electron gas.

At a high concentration of donor mixtures, the expanding penetration zone may overlap with the conductivity zone (or with the tensile zone in the semiconductor's capacitor). Thus, a single zone of allowable energy values is formed in the crystal.

The presence of an ingress zone affects the density of states in the permeability zone close to E_c (or close to E_v fa, respectively). At the ordinarily ingress concentration, the bandit coating is small, and the energy dependence of the state density close to the 10^{19} sm^{-3} E_c (or E_v) level is significantly distorted, and the conductivity band (or tensile zone) becomes non-standard, as shown in Figure 4.20. In such a semiconductor, due to high-order legging, a "tail" of state density is formed, and the exclusion zone E_g is reduced by E_g . In a circuit where there is a

degree of legering, the narrowing of the pronouncement zone is a few hundredths of a volt of an electron. As the level of leprosy increases, the effect of the "tailing" of state density increases. This 10^{19} sm^{-3} follows from Fig. 4.21, which presents the results of a study to determine the density of cases in p-type gallium arsenide obtained by tunneling in a metal-semiconductor contact.

Despite the fact that fused semiconductors have very wide applications in some semiconductor devices (e.g., tunneling diodes, quantum generators), the rigorous theory of strongly alloyed semiconductors is a mathematically difficult task. Experiments have shown that in n or p-type semiconductors, electrical conductivity at low temperatures can also be observed at the concentration of the mixture. This mechanism of electrical conductivity is called jumping. Let's look at this in the case of an electron semiconductor. A true donor semiconductor always has a certain number of $10^{16} \text{ sm}^{-3} N_d$ receptor centers, so that some electrons pass from the N_d donor input layer to the receiver level. As a result of such compensation, even at very low temperatures, donors in the amount of N_d turn out to be ionized. Therefore, an electron localized in one of the donors can "jump" to the nearest ionized donor, thereby ionizing the former and neutralizing the latter. This circuit jump happens randomly. When an external electric field is present, the jumping of electrons becomes directional, resulting in a charge current. Since the donors are located at large distances from each other, the mobility of charge carriers driven by the jumping mechanism is significantly less than the mobility of electrons in the conduction zone. In the area of high temperatures, normal permeability prevails across energy zones. Therefore, over the entire temperature range the permeability of such a sample σ and the Hall coefficient R are determined by expressions such as:

$$\sigma = en_1\mu_1 + en_2\mu_2 \quad (4.143)$$

$$R = \frac{r_1 n_1 \mu_1^2 + r_2 n_2 \mu_2^2}{e(n_1 \mu_1 + n_2 \mu_2)^2} = \frac{n_1 \mu_1 \mu_{1H} + n_2 \mu_2 \mu_{2H}}{e(n_1 \mu_1 + n_2 \mu_2)^2}, \quad (4.144)$$

Where $r_1 = \frac{\mu_{1H}}{\mu_1}$; $r_2 = \frac{\mu_{2H}}{\mu_2}$, a, $n_1, \mu_1 \mu_{1H}$ and is the concentration in the conduction zone and the input conductivity region, drift and Hall mobility. n_2, μ_2, μ_{2H}

4.10. Noncrystalline semiconductors

In non-crystalline amorphous semiconductors, only the short-distance order is preserved, so that within several coordinate spheres, the arrangement of atoms in crystalline and amorphous material is almost the same. But as one moves away from the atom taken as a starting point, the effect of the random factor in the amorphous

material becomes more intense, and it is impossible to establish a rigorous geometric relationship between two atoms at a great distance from each other, which is characteristic of crystals. In non-crystalline materials, the atomic environment varies randomly from atom to atom, and therefore they are characterized by the absence of long-range order in the arrangement of atoms. Therefore, the concept of mean free path cannot be used for amorphous material, and the concept of Brillouen zones is also not applicable. However, it is theoretically shown that in non-crystalline materials, there may be a state of electrons in which the wave functions extend into the macroscopic dimensions domain. These states, called "extended" states, can form a band of energy in a crystal. In addition, there are bonded states in amorphous materials where the dimensions of the wave function are localized in a region where the size of the wave function is not much greater than the interatomic distance. The energy levels corresponding to the localized states are located near the edges. In Fig. 4.22, which shows the density of the $N(E)$ states, the energies E_c and E_v coincide with the edges of the crystalline material, and the fixed state of E_t and E_{tv} . The areas occupied by localized cases are marked by bars. In addition, the density of cases can cross a large number of maximums that are significantly far from each other (Fig. 4.22- b). For the $N(E)$ distribution shown in Fig. 4.22, the "tails" of the bands stick together in such a way that the entire interval between E_{tc} and E_{tv} contains energy levels of localized states. Since electrical conductivity is prohibited in localized states, the mobility of charge carriers in these cases is zero. Therefore, the term "forbidden zone" for amorphous materials does not correspond to the energy structure of the layers, and it is common to speak of a "void" mobility zone, the value of which (Fig. 4.22-d) exceeds the width of the forbidden zone in the crystalline material. $\Delta E_\mu = E_{tc} - E_{tv} = E_c - E_v$

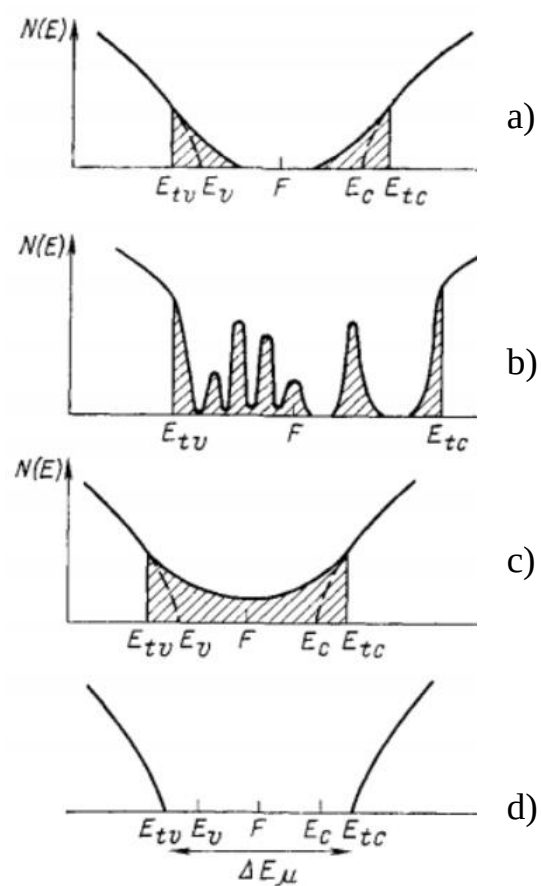


Fig. 4.22. Zone changes in amorphous semiconductor: a-zone tails; b-The possible distribution of the density of the cases; intersection of V-tails; G-man's zone.

Table 4.1

Close Atom Count	Group in periodicals	Material
------------------	----------------------	----------

2	A^{VI}	Se, Te
3	A^V	Ace, Sb, P
3 & 4	A^{IV}	
4	A^{IV}	Ge, Si
6	A^{III}	B
3-2	$A^{d \ b \ VI}$	$(As, Sb, Bi)_2(S, Se, Te)_3$
4-2	$A^{IV} - B^{VI}$	$SiO_2, (Ge, Sn, Pb)(S, Se, Te)$

Experimental work in the field of amorphous semiconductors has revealed a close relationship between the electrical and optical properties of these materials and the distribution of localized cases, as well as a close relationship between the density of cases in the bandit mobility zone and the material's fabrication technology. An example of the impact of material preparation technology can be seen in Fig. 4.23 with the gloss resorption method (Curve 1) and the vacuum immersion method (Curve 2) for amorphous silicon (α -Si) films.

It follows from the graphs that when the technology of producing α -Si films is changed, the intensity of localized states in the mobility gap can change by several degrees. These states are neutral when they are not filled with electrons. Below the Fermi level, these states trap electrons and are negatively charged. Above the Fermi level, these cases are positively charged. For most amorphous materials, the rule of thumb is that the first coordinate number of the corresponding crystal is retained in the amorphous material. This basic rule can be used to classify amorphous materials. Table 4.1 shows the distribution of a number of amorphous materials according to the number of coordinates.

As can be seen from the data in the table, amorphous materials, even in the first coordination area, can have great structural diversity. And yet, according to their structural characteristics, they can be divided into two classes. These are materials with tetrahedral coordination and chalcogenide glass type. The materials of the first group are characterized by a rigid structural framework with a very limited capacity for the reorganization of local atoms.

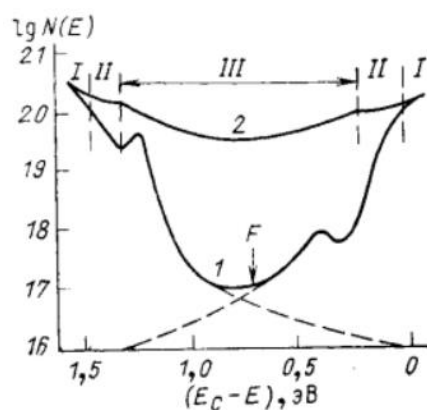


Fig. 4.23. State density distribution in amorphous films obtained by spark refraction and spraying

Materials of the second group, as a rule, are characterized by a smaller number of the nearest neighbors and have a greater elasticity. Tetrahedral material of type a-Si is characterized by a large number of hanging links (commonly called "hanging" bonds). The presence of unpaired spin electrons leads to the release of the EPR signal, which gives a concentration of 10^{19} - 10^{20} cm^{-3} for the density of the paramagnetic centers. The introduction of hydrogen into these materials leads to saturation of the hydrogen bonds and loss of the EPR signal. This material can be used with donor and acceptor mixtures and it is used to manufacture some semiconductor devices. Materials of the second group (e.g., selenium and chalcogenide pores) do not emit an EPR signal, indicating the absence of paramagnetic centers, or at least their density below the detectable limit of 10^{15} cm^{-3} . In this case, the existing suspended bonds must be filled with two electrons with an antiparallel spin.

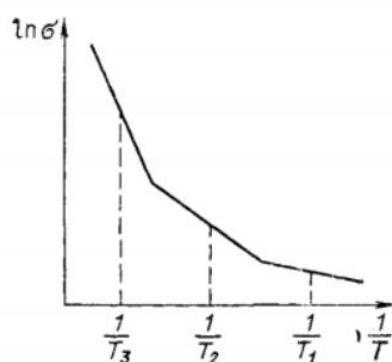


Fig. 4.24. Temperature dependence of conductivity in amorphous materials

Another feature of chalcogenide materials is that they have six electrons in their outer shell, and not all of them are involved in bond formation. For example, for Se, two electrons are involved in the formation of bonds, four electrons are in the unbonded state, i.e., the so-called single pairs. These are the electrons of the

unallocated pairs that make up the stress zone of chalcogenide glass semiconductors. The presence of solitary pairs determines the energy structure and physical properties of glassy materials. In amorphous semiconductors, at temperatures other than zero, electrical conductivity can be achieved via a jump mechanism through localized states. Since the probability of the transition of electrons from one atom to another depends on temperature, this conductivity is thermally activated (Fig. 4.24).

The nature of temperature dependency is determined by the hopping of an electron at a given temperature and what group of qualms its mechanism depends on. The optical properties of amorphous samples depend more on the conditions of their preparation than on the properties of crystals. This is explained not only by the fact that an amorphous material can significantly change its density curve with a change in preparation conditions, but also the appearance of micropores (especially in tetrahedral materials), which can cause scattering of light in the sample. In glassglass semiconductors, the effect of current stability is observed, on the basis of which it is possible to create switches and memory devices with a very short changeover time 10^{-10} s.

Control Questions

1. What is the density of quantum states?
2. Explain the Fermi-Dirac distribution?
3. Explain the Boze-Einstein distribution?
4. How is the concentration of electrons and holes in the zones determined?
5. How does the Fermi level relate to the concentration of inputs?
6. How is the Fermi level related to the temperature of an unchanged semiconductor?
7. Give examples of non-crystalline semiconductors.

Chapter V

SCATTERING OF ELECTRONS AND HIVES IN SEMICONDUCTORS

5.1. Electrons and Scattering Mechanisms of Spurs

When considering the motion of charge carriers in the ideal circuitous field of a crystal lattice, it was found that electrons and spurs act as free particles (see Fig. 2.10). In a true crystal, electrons and spikes go through complex trajectories due to collisions with lattice defects. Since defects disrupting the periodicity of the lattice area and having scattering centers are of a different nature, they also give rise to different mechanisms of dissipation of charge carriers. In semiconductors, scattering centers are the thermal vibrations of the lattice and static defects such as input atoms and ions, voids, dislocations, twin boundaries, and crystallites. In order to quantitatively evaluate the sowing process, a parameter called effective shaving section σ is introduced.

Suppose there are n free electrons moving in a certain direction with an average heat velocity v_0 . Then nv_0 is the density of the electron current, that is, the number of electrons that pass through one area of the sample perpendicular to the direction of their velocity in a unit of time. Suppose that in one section of the sample there are N identical scattering centers along the path of the electron current. Each center σ is characterized by an equally effective cross-section. This is, by its essence, the space around the center where the scattering of electrons occurs. Therefore, the number of scattered electrons in the unit of time n_1 is determined by the effective cross-section σ , the number of scattering centers N , and the density of the incident electron current nv_0 , i.e.

$$n_1 = \sigma N n v_0 \quad (5.1)$$

If W is the probability of scattering of a single particle in unit of time, then the number of electrons scattered in unit of time W

$$n_1 = W n \quad (5.2)$$

is equal to n_1 . Then, according to (5.1) and (5.2)

$$\sigma = \frac{n_1/N}{n v_0} = \frac{W}{N v_0}. \quad (5.3)$$

It can be written.

Thereby, the effective scattering section is the ratio of the σ number of electrons released from the current by a centrifugal scattering in unit of time to the density of the falling particle stream. The effective scattering section has the dimensions of the area

$$[\sigma] = \frac{[W]}{[Nv_0]} = \frac{T^{-1}}{L^{-3}LT^{-1}} = L^2.$$

Probability of scattering from (5.3)

$$W = \sigma N v_0. \quad (5.4)$$

is equal to . Consequently, the probability of scattering or collision is determined by the effective cross-section, the number of scattering centers, and the speed of movement of the zaaryad carriers.

At the same time, the probability of a collision is inversely proportional to the time of a free run

$$W = 1/\tau. \quad (5.5)$$

Therefore, the free-running time and free-running path can be expressed through an effective cross-section. (5.4) and (5.5)

$$\tau = \frac{1}{\sigma N v_0} = \frac{l}{v_0} \quad (5.6)$$

or

$$l = \frac{1}{\sigma N} \quad (5.7)$$

origin. probability of scattering in the magnitude unit path interval. $l^{-1} = \sigma N$

Suppose that there are different scattering centers. Let the i -center be characterized by an effective cross-section σ_i and let the number of such centers be N_i . The swing mechanism in these centers determines the average free-running path. If all scattering processes are possible, according to probability theory, the total probability of scattering in a unit of time is determined by the sum of the individual scattering probabilities.

$$W = \sum_i W_i. \quad (5.8)$$

$\sum_i \sigma_i N_i = \sum_i l_i^{-1}$ Based on (5.4) and (5.8) free runs.

$$l^{-1} = \sum_i l_i^{-1}. \quad (5.9)$$

can be determined by attitude. Based on (5.9), the full free-running path will always be smaller than the partial free-running track.

Since the role of defects in the process of charge carrier scattering is different, different defects must have different effective cross-sections. To quantify them, we take the effective scattering section as the area where the interaction between the charge carrier and the defect is possible.

Defects such as vacancies and internode atoms are in many ways like substitute mixtures. These defects are called point defects. For them, σ_i can be taken as the area of a square whose side is equal to the grid constant, i.e., $\sigma_A =$

$(5 \cdot 10^{-8})^2 \approx 3 \cdot 10^{-15} \text{ sm}^2$. If we assume the concentration of the inlet atoms, the mean free path of the charge carriers scattered into the ingress atoms is as follows: $N_A = 10^{16} \text{ sm}^{-3} l_A = (3 \cdot 10^{-15} \cdot 10^{-16})^{-1} \approx 3 \cdot 10^{-2} \text{ cm} = 300 \text{ } \mu\text{m}$.

We assume that the diameter of the inlet ion is 10 times greater than the diameter of the inlet center, that is, . If it is, it will be mkm . $\sigma_I = (5 \cdot 10^{-8} \cdot 10)^2 \approx 3 \cdot 10^{-13} \text{ cm}^2 N_I = 10^{16} \text{ cm}^{-3} l_I = 3 \cdot 10^{-4} \text{ cm} = 3$

Dislocations are linear defects that extend over a large area of the crystal. Suppose that a linear dislocation has a length of 0.1 cm and its diameter is equal to the constant of 100 grids. In this case, its cross-sectional surface is equal to the surface. At the volumetric density of the dislocations, the free-running path is $\text{cm } \mu\text{m} \cdot 5 \cdot 10^{-8} \cdot 100 \cdot 10^{-1} = 5 \cdot 10^{-7} \text{ cm}^2 N_D = 10^8 \text{ cm}^{-3} l_D = 2 \cdot 10^{-2} = 200$

The effective cross-section of scattering in the thermal oscillations of the lattice is determined by the subsection of the area occupied by the oscillation atom from the cross-section of the area occupied by the atom itself. If the amplitude of the vibrations $r = 0,05$ is nm cm , the atomic diameter is cm , then cm . This value is much smaller than the value of the other scattering centers of scattering, but the number of atoms oscillating is great: , so $\mu\text{m} = 5 \cdot 10^{-9} d = 10^{-8} \sigma_T = (d + r)^2 - d^2 \approx 2rd = 10^{-16} \text{ }^2 N_T = 10^{22} \text{ cm}^{-3} l_T = 10^{-6} \text{ cm} = 0,01$

5.2. Boltzmann's kinetic equation

Until now, the set of charge carriers in a crystal has been considered as a system of thermodynamically balanced particles. In this case, the motion of the charge carriers is chaotic, and their behavior is characterized by a balanced distribution function, which is usually dependent on the energy of the particles:

$$f_0 = \frac{1}{e^{(E-F)/kT} + 1} \quad (5.10)$$

If the crystal is placed in an external constant electric field, the movement of the charge carriers will be orderly.

Phenomena that occur as a result of the movement of charge carriers under the influence of electric and magnetic fields, temperature gradient and concentration gradient of charge carriers are called migration phenomena or kinetic phenomena. Landshear phenomena include electrical conductivity, electron thermal conductivity, galvanomagnetic, thermoelectric, and thermomagnetic effects.

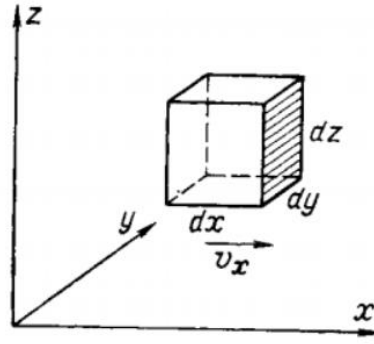


Fig. 5.1. Elementary volume element.

In the presence of external forces, the charge carriers are in a state of disproportion. Naturally, in the unbalanced state, the charge carrier system is characterized by an unbalanced distribution function, which must depend on the energy of the particles according to (2.102) and Figure 2.13, i.e., k and their coordinates r and time t , so the function $f(k, r, t)$ must be dependent.

To determine the concentration of electrons in the non-equilibrium state, we proceed as follows. We select the volume element dG in the spatial phase of the unit-volume crystal, i.e. the space of x, y, z coordinates and quasi-impulses. This p_x, p_y, p_z would be a space built around a point with p_x, p_y, p_z coordinates (x, y, z) :

$$dG = dx dy dz dp_x dp_y dp_z = d\tau_r d\tau_p = \hbar^3 d\tau_k d\tau_r \quad (5.11)$$

here

$$d\tau_r = dx dy dz; d\tau_p = dp_x dp_y dp_z = \hbar^3 dk_x dk_y dk_z = \hbar^3 d\tau_k$$

In the volumetric element of the phase phase of a unit-volume crystal, there are $dG/\hbar^3$ phase cells, each of which may contain two electrons with an antiparallel spin. Therefore, the volume element dG of the phase phase contains the quantum state of $2dG/\hbar^3$. If $f(k, r, t)$ is the probability of finding an electron in these states, then the number of electrons in an element of phase space at time t will be equal to:

$$dn = 2 \frac{d\Gamma}{\hbar^3} f(\mathbf{k}, \mathbf{r}, t) = \frac{1}{4\pi^3} f(\mathbf{k}, \mathbf{r}, t) d\tau_k d\tau_r. \quad (5.12)$$

(5.12) We determine the total number of electrons by integrating the expression across the entire space

$$n = \frac{1}{4\pi^3} \int_{(V_r)} f(\mathbf{k}, \mathbf{r}, t) d\tau_k d\tau_r. \quad (5.13)$$

Thus, when considering kinetic phenomena, it is necessary to know the unbalanced distribution function. We derive an equation satisfying the unbalanced distribution function $f(k, r, t)$. To do this, we consider the motion of electrons in ordinary space and the space of wave factors.

First, let's consider a system of electrons in a simple geometric space. Suppose the external electric field is applied in such a way that electrons move in the *positive* direction of the x -axis with a velocity v_x . We determine the change

in the concentration of electrons within the volume element D in time dT . According $\tau_{to(5.2)}$ shown in Fig. 5.1, the volume element D r includes electrons moving with a VX velocity through the left side of the time dt :

$$dN_{\pi} = f(\mathbf{k}, x, y, z, t) \frac{d\tau_k}{4\pi^3} v_x dydzdt \quad (5.14)$$

During this time, the change in the number due to the departure of electrons from the right boundary of the pilgrimage

$$dN_{II} = -f(\mathbf{k}, x + dx, y, z, t) \frac{d\tau_k}{4\pi^3} v_x dydzdt \quad (5.15)$$

$d\tau_r$ Change in the number of electrons in a volume element over time d

$$dN_n + dN_n = [f(\mathbf{k}, x, y, z, t) - f(\mathbf{k}, x + dx, y, z, t)] \frac{d\tau_k}{4\pi^3} v_x dydzdt =$$

In general, the change in the number of electrons in the unit of time in the volume element dg is the result of the movement of charge carriers with a wave vector k in the arbitrary direction with velocity v

$$-\left[v_x \frac{\partial f}{\partial x} + v_y \frac{\partial f}{\partial y} + v_z \frac{\partial f}{\partial z}\right] \frac{d\tau_k}{4\pi^3} d\tau_r dt = -(\mathbf{v}_{\nabla_r} f) \frac{d\tau_k}{4\pi^3} d\tau_r dt \quad (5.17)$$

is defined by equality. Such a change in the number of electrons is due to the displacement phenomena that occur by diffusion of charge carriers at the expense of concentration or temperature gradient

Let us consider the motion of electrons in the space of wave vectors. Let us calculate that the number of electrons varies in time dt τ due to the arrival and departure of the volume element dk through the appropriate facets of the volume element. To do this, just as in the motion of electrons in ordinary space, we need to calculate the number of electrons that come through and remain through all the faces of the volume element $d\tau_k$. As a result, when electrons move in the wave vector space, the number of volumes dG in the element of the phase space in dt time changes to the following value:

$$-\left[\frac{\partial k_x}{\partial t} \frac{\partial f}{\partial k_x} + \frac{\partial k_y}{\partial t} \frac{\partial f}{\partial k_y} + \frac{\partial k_z}{\partial t} \frac{\partial f}{\partial k_z}\right] \frac{d\tau_k}{4\pi^3} d\tau_r dt =$$

here

$$\mathbf{F}(\mathbf{r}, t) = \frac{d\mathbf{p}}{dt} = \hbar \frac{d\mathbf{k}}{dt} \quad (5.19)$$

- An rt external force acting on an electron at a moment of time at a point.

The relation (5.18) represents the change in the number of electrons as a result of the action of external forces caused by the action of electric and magnetic fields.

In addition, the number of electrons in the volume element of the phase space can be varied by dG because of the interaction of electrons with local disturbances

in the periodicity of the crystal lattice area. It can be assumed that the range of influence of local disturbance extends at a distance of about 10^{-7} cm, covering several lattice periods. Since the thermal velocity of the electron in the semiconductor at room temperature is equal to 10^7 cm/s, the interaction time of an electron with a structure defect in the crystal is to be 10^{-14} s. Such a short-term interaction is unlikely to cause a significant change in the coordinate, but it will significantly change the velocity and quasi-impulse of the electron. This process is similar to a shock in mechanics, so it is called a collision or shock. During collision, a change in the number of electrons moving in a given direction occurs, so collisions are also called scattering processes.

Suppose that during decay, electrons go from (r, k) states to (r', k') states. Since the collision involves a short-term interaction that does not cause a change in the coordinate, the electron jump probability will not depend on r and r' . We define the probability of such passage in a unit of time as $W(k, k')$. The number of states occupied in the element $d\mathbf{k}$ is $f(k, r, t)d\tau_k/4\pi^3$, the number of free states in this volume is respectively $[1 - f(k, r, t)]d\tau_k/4\pi^3$. Each collision is accompanied by the transition of the electron from the k state to the k' state. But for such a transition, the state *must be* k' free. If $f(k', r, t)d\tau_{k'}/4\pi^3$ is the number of cases occupied by the element $d\tau_{k'}$, then $[1 - f(k', r, t)]d\tau_{k'}/4\pi^3$ is the number of empty states. Therefore, the number of electrons capable of switching from state k to state k' as a result of dissipation in time dt leads to a decrease in the number of electrons in the phase space volume element dG , the transition probability is determined by $W(k, k')$, the number of initial states occupied, and the number of free final states:

$$-W(\mathbf{k}, \mathbf{k}')f(\mathbf{k}, \mathbf{r}, t)\frac{d\tau_{\mathbf{k}}}{4\pi^3}[1 - f(\mathbf{k}', \mathbf{r}, t)]\frac{d\tau_{\mathbf{k}'}}{4\pi^3}d\tau_{\mathbf{r}}dt \quad (5.20)$$

It is also possible for electrons to switch in reverse from the k' state to the k state when the charge carriers collide with the lattice defects. Let the probability of such a transition in a unit of time be $W(k', k)$. Then the number of electrons that have completed the reverse transition in time dt is equal to:

$$W(k', k)f(k', r, t)\frac{d\tau_{k'}}{4\pi^3}[1 - f(k, r, t)]\frac{d\tau_k}{4\pi^3}d\tau_rdt \quad (5.21)$$

Therefore, the number of electrons in the element dg changes as a result of scattering, along with direct transitions from k to k' and vice versa from k to k . If f is expressed as a function of the wave vector to shorten the record, then dt is the change in the number of electrons over time

$$\{W(k', k)f(k')[1 - f(k)] - W(k, k')f(k)[1 - f(k')]\}\frac{d\tau_{k'}}{4\pi^3}\frac{d\tau_k}{4\pi^3}d\tau_rdt \quad (5.22)$$

So so.

As a result of scattering, all possible k must be taken into account in order to find the total number of states occupied by electrons in the volume element of the phase space in dt time, i.e., the expression in terms of the size of the Brillouin zone (5.22) must be integrated.

$$\frac{d\tau_k}{4\pi^3} d\tau_r dt \int_{V_B} \{W(k',k) f(k')[1-f(k)] - W(k,k') f(k)[1-f(k')]\} \frac{d\tau_{k'}}{4\pi^3} \quad (5.23)$$

But the change in the number of electrons in the volume element dG of the phase space in time from t to $t+dt$

$$f(k,r,t+dt) \frac{d\tau_k}{4\pi^3} d\tau_r - f(k,r,t) \frac{d\tau_k}{4\pi^3} d\tau_r dt = \frac{\partial f}{\partial t} \frac{d\tau_k}{4\pi^3} d\tau_r \quad (5.24)$$

Such a variation in the number of charge carriers is due to the movement of electrons due to diffusion, the effect of an external force on them and the presence of scattering processes. Thus, the expression (5.24) is equal to the sum of the expressions (5.17), (5.18), and (5.23). After the abbreviation, we get:

$$\frac{\partial f}{\partial t} = -(\mathbf{v} \nabla_r f) - \frac{1}{\hbar} (F \nabla_k f) + \int_{V_B} \{W(k',k) f(k')[1-f(k)] - W(k,k') f(k)[1-f(k')]\} \frac{d\tau_{k'}}{4\pi^3} \quad (5.25)$$

Based on the principle of the reversibility of forward and reverse transitions

$$W(k,k') = W(k',k). \quad (5.26)$$

(5.26) vs. (5.25)

$$\frac{\partial f}{\partial t} = -(\mathbf{v} \nabla_r f) - \frac{1}{\hbar} (F \nabla_k f) + \int_{V_B} W(k',k) f(k')[f(k) - f(k')] \frac{d\tau_{k'}}{4\pi^3} \quad (5.27)$$

takes the view.

The equation (5.27) is commonly referred to as the Boltzmann kinetic equation. This equation is integro-differential because the desired function is under the integral.

It is easy to see the physical meaning of the equation. Its left side reflects the change in the distribution function over time. The first and second terms on the right side of the equation show the change in the distribution function that leads to the motion of electrons in normal space and in the space of wave vectors under the influence of external fields. The change in the distribution function across states as a result of scattering (collision) of electrons is considered as the last case in the equation. Thus, charge carriers are affected on the one hand by external forces, which tend to regulate their motion, and on the other, collisions, which tend to disrupt their motion. Therefore, the equation (5.27) can be written as

$$\frac{\partial f}{\partial t} = \left(\frac{\partial f}{\partial t} \right)_{\text{coll}} + \left(\frac{\partial f}{\partial t} \right)_{\text{st}} \quad (5.28)$$

here

$$\left(\frac{\partial f}{\partial t}\right)_{to'la} = -(\mathbf{v}\nabla_r f) - \frac{1}{\hbar}(F\nabla_k f) \quad (5.29)$$

It is called the field organizer of the Boltzman equation.

$$\left(\frac{\partial f}{\partial t}\right)_{st} = \int_{V_B} \{W(k, k')[1f(k') - f(k)]\frac{d\tau_{k'}}{4\pi^3}\} \quad (5.30)$$

And the magnitude is called the collision integral.

In stationary processes, when quantities describing displacement events are not time-dependent, the distribution function

$$\frac{\partial f}{\partial t} = 0 \quad (5.31)$$

satisfies PCB and Boltzman's for stationary status (5.28)

$$\left(\frac{\partial f}{\partial t}\right)_{to'la} = -\left(\frac{\partial f}{\partial t}\right)_{st} \quad (5.32)$$

you can change the look, or

$$(\mathbf{v}\nabla_r f) + \frac{1}{\hbar}(F\nabla_k f) = \int_{V_B} W(k', k) f(k')[f(k') - f(k)]\frac{d\tau_{k'}}{4\pi^3} \quad (5.33)$$

takes the view.

It follows from the expressions (5.32) and (5.33) that the changes in the unbalanced distribution function caused by the motion of the charge carriers in the stationary state and the impact of external force on them are compensated by the local violation of the circuit of the mesh field with the collision of the charge carriers.

It should be noted that the Boltzman kinetic equation is essentially quasi-elastic, since the coordinate and the wave vector are simultaneously defined in the distribution function. Therefore, it applies to processes in which the concept of a moderate free path makes sense. This means that the average free path of the charge carriers must be greater than the de Broyle wavelength. In addition, this equation can be used to describe the phenomenon of a transition in a magnetic field, provided that this field does not change the energy spectrum of the charge carriers in the crystal, i.e., that the quantum effects are not significant. This condition places constraints on the value of the magnetic field strength.

5.3. Equilibrium state

In thermodynamic equilibrium, where the state of electrons is characterized by the equilibrium distribution function (5.10), the movement of electrons in a crystal occurs under the action of heat, so for the equilibrium state (5.33) the left side of Equation (5.33) is zero. This means that

$$\int_{V_B} W(k', k) [f_0(k') - f_0(k)] d\tau_{k'} = 0 \quad (5.34)$$

It follows from this relationship that in the system in the equilibrium state there is an equivalence of charge carrier currents for the direct and inverse processes, which corresponds to the principle of equilibrium in detail.

Based on (5.34), we get:

$$f_0(k') = f_0(k) \quad (5.35)$$

Balanced distribution function by subtracting (5.10) by (5.35),

$$E_1 - F_1 = E_2 - F_2, \quad (5.36)$$

We have E_2 , where E_2 , $-k$ and k' are the full energy for electrons with the wave vector and the Fermi level. Under thermodynamic conditions, the total energy of an electron does not change $E_1 = E_2$, so from (5.36)

$$F_1 = F_2 \quad (5.37)$$

origin.

This expression indicates that in the equilibrium state, the Fermi level will be the same in all parts of the system where electrons can pass.

5.4. Relaxation time

As we know, the solution of the kinetic equation (5.27) is very difficult even for the stationary state (5.33). But this simplifies for situations where a relaxation time can be inserted. It turns out that if the collision of particles is elastic, that is, if the energy of the charge carrier is slightly changed, and the scattering processes result in a distribution of the charge carriers by velocities, that is, equally probable across all directions of the charges, then it is possible to describe the scattering processes using the relaxation time.

Suppose that the function of the unbalanced distribution f at the moment of time $t = 0$ f , the external distortions cease to act (zeros are erased) and the field becomes zero. As a result of collisions, the system arrives at an equilibrium state, which is characterized by the distribution function f_0 . This means that after the outer field is removed, the distribution function changes due to the presence of collisions of electrons with lattice defects:

$$\frac{\partial f}{\partial t} = \left(\frac{\partial f}{\partial t} \right)_{st} \quad (5.38)$$

If the deviation of the charge carrier distribution from the equilibrium state is small, in the absence of external fields, we can assume that the rate of change of the distribution function due to collisions is proportional to the magnitude of the deviation from equilibrium of the function, i.e., proportional to f :

$$\frac{\partial f}{\partial t} = \left(\frac{\partial f}{\partial t}\right)_{st} = \frac{f - f_0}{\tau(k)} \quad (5.39)$$

Here is the proportionality coefficient dependent on $\tau(k)$.

Solving the equation 5.39 we get:

$$f - f_0 = (f - f_0)_{t=0} e^{-t/\tau} \quad (5.40)$$

It follows from (5.40) that once the influence of external fields is stopped, the difference $(f - f_0)$ decreases according to the indicative law with a time constant called relaxation time τ . Consequently, τ it is the average time that there is an unbalanced distribution of charge carriers in the system after the outer fields have been removed. Since the establishment of the equilibrium state is due to collisions and only a few collisions are sufficient, the relaxation time must be approximately equal to the free travel time of the circuit. Taking this into account, it can be assumed that τ is determined by the mean free-path electron velocity v and the relaxation time (k) .

Using (5.39) and (5.30) equation can be:

$$\left(\frac{\partial f}{\partial t}\right)_{st} = \int_{V_B} W(k, k') [f(k') - f(k)] \frac{d\tau_{k'}}{4\pi^3} = \frac{f_0(k) - f(k)}{\tau(k)} \quad (5.41)$$

In accordance with the equation (5.39), since the external forces do not cause large deviations of the charge carrier distribution from the equilibrium state, the unbalanced distribution function can be expressed as follows.

$$f(k) = f_0 + f_1(k), \quad (5.42)$$

where $f_1(k)$ - is a correction to the equilibrium distribution function resulting from an external influence. We write it as follows

$$f_1(k) = -\frac{\partial f_0}{\partial E} k \chi(E) \quad (5.43)$$

Here $\chi(E)$ - is some kind of vector that depends on the energy of the charge carrier.

We assume that the relaxation time does not depend on the outside areas. Hence, $\tau(k)$ also determines the scattering processes involving external fields. Therefore, for equation (5.33), we can use the collision integral of (5.41), and for the stationary state the kinetic equation takes the following form:

$$(v \nabla_r f) + \frac{1}{\hbar} (F \nabla_k f) = \frac{f_0 - f}{\tau(k)} = \frac{f(k)}{\tau(k)} \quad (5.44)$$

Using the relation (5.42), we write the collision interval as follows

$$\left(\frac{\partial f}{\partial t}\right)_{st} = \frac{1}{4\pi^3} \int_{V_B} W(k, k') [f_1(k') - f_1(k)] d\tau_{k'} = -\frac{f_1(k)}{\tau(k)} \quad (5.45)$$

(5.45) we can determine the time of relaxation from the equation. If (5.43) takes into account the attitude, (χk) is determined by equality of the following form:

$$\frac{1}{\tau(k)} = \frac{1}{4\pi^3} \int_{V_B} W(k, k') \left[1 - \frac{\partial f_0 / \partial E, k' \chi(E')}{\partial f_0 / \partial E, k \chi(E)} \right] d\tau_{k'} \quad (5.46)$$

As we already know, collisions can be elastic and inelastic. If the kinetic energy of the intercolliding particles is conserved, the collision is called elastic. Noelastic collisions are collisions in which the energy of the particles changes and is more or less the same as before the impact after the impact.

With structural distortions, we assume that the collision of electrons is elastic. In this case, only the direction of the wave vector and therefore the direction of the velocity of the electron changes, and the absolute values of their motion remain unchanged. Therefore, during the elastic collision

$$k' = k; \quad v' = v; \quad E' = E, \quad (5.47)$$

of the Enlightenment Project. Therefore, the relaxation time for elastic collision given (5.47)

$$\frac{1}{\tau(k)} = \frac{1}{4\pi^3} \int_{V_B} W(k, k') \left[1 - \frac{k' \chi(E')}{k \chi(E)} \right] d\tau_{k'} \quad (5.48)$$

will be equal to .

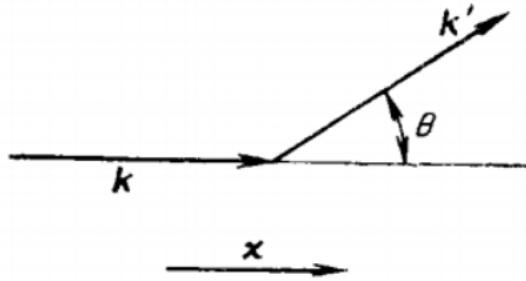


Fig. 5.2. Change of an electron wave vector in scattering

Let the vectors k and k' be oriented as shown in Fig. 5.2. In this case, the projection of vectors k and k' onto the vector x will be,

$$k_x = k; \quad k'_x = k' \cos \theta = k \cos \theta \quad (5.49)$$

where θ - is the deviation of the electron from the initial direction. (5.49) parity given (5.48) is written in the form:

$$\frac{1}{\tau(k)} = \frac{1}{4\pi^3} \int_{V_B} W(k, k') \left[1 - \frac{k'}{k} \right] d\tau_{k'} = \int_{V_B} W(k, k') [1 - \cos \theta] d\tau_{k'} \quad (5.50)$$

Thus, the main task in considering the processes of transfer is to calculate the unbalanced distribution function, with the help of which it is possible to calculate all quantities describing kinetic effects. However, from (5.44) it is possible to solve this problem when the relaxation time (5.50) is known, and therefore it is necessary to know the dissipation mechanism of the charge carrier.

Let's consider scattering, which causes a random distribution of charge carriers across velocities, i.e., where the scattering of an electron ends in the state k' , which is close to the same energy minimum as k , but with an inversely directed velocity. Such scattering is possible if the surfaces of constant energy are spherical.

Since $W(k, k') = W(k', k)$ - is a double function, $f_1(k') = -\partial F_0 / \partial (\hbar k' \chi)$ - is an odd function

$$\int_{V_B} W(k, k') f_1(k') d\tau_{k'} = - \int_{V_B} \frac{\partial f_0}{\partial E} W(k, k') k' d\tau_{k'} = 0 \quad (5.51)$$

(5.51) Expression given (5.45)

$$\left(\frac{\partial f}{\partial t} \right)_{st} = - \frac{1}{4\pi^3} \int_{V_B} W(k, k') f_1(k') d\tau_{k'} = - \frac{\partial f_1(k)}{\tau(k)} \quad (5.52)$$

takes the view because when integrated within symmetric boundaries, the odd function returns zero. It follows that the relaxation time during a scatter, which causes a random distribution of the charge carriers by velocities is defined by expression.

$$\frac{1}{\tau(k)} = \frac{1}{4\pi^3} \int_{V_B} W(k, k') d\tau_{k'} \quad (5.53)$$

5.5. Scattering in Input Ions

In an inlet semiconductor, each input ion generates an electric field around it. Under the influence of this field, the moving charge carrier deviates from its original orientation (Fig. 5.3). In this case, the slower and closer it moves to the ion, the stronger the charge carrier will deviate.

The problem is actually finding the trajectory of a charged particle in the central area of the ion. This can be solved only with the help of classical methods. To do this, we consider the charge carrier to be a classical particle moving with a Ze charge in the Coulomb field of the positive input ion. In this case, the potential energy of the interaction

$$U(r) = \pm \frac{Ze^2}{4\pi\epsilon_r\epsilon_0 r} \quad (5.54)$$

can be expressed in the form, where plus is the hole, minus the electron.

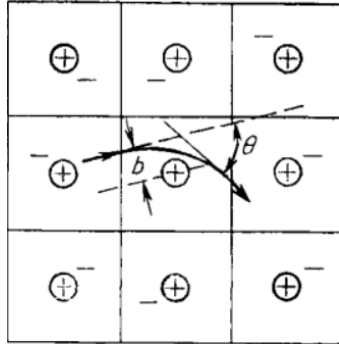


Fig. 5.3. Curvature of the electron motion trajectory in the positive input ion field

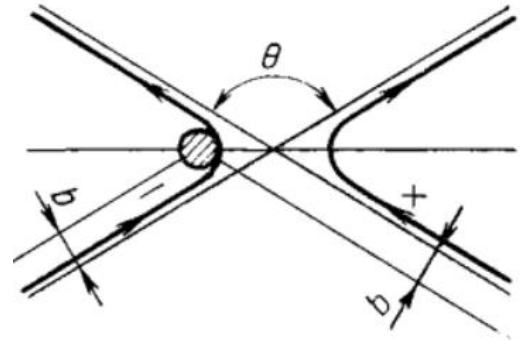


Fig. 5.4. Electron and Spark Scattering in Positive Input Ion

Calculations show that in the Coulomb field of the ion, the charge carrier moves along the hyperbola. Fig. 5.4 shows the trajectories of the electron and the spark's motion in the positive ion field. Both charge carriers bend equally, although in one case there is a pulling force and in the other there is a repulsive force, b refers to the striking distance - the minimum distance between the ion and the trajectory of the charge carrier, and θ - the angle of scattering, i.e., the angle between the initial direction of motion of the charge carrier and the direction of motion after scattering. In this case, the angle of sweep depends on the distance as follows:

$$b = \pm \frac{Ze^2}{\epsilon_r m^* v^2} \cotg \frac{\theta}{2} \quad (5.55)$$

Here t^* and v - are the effective mass and velocity of the charge carrier.

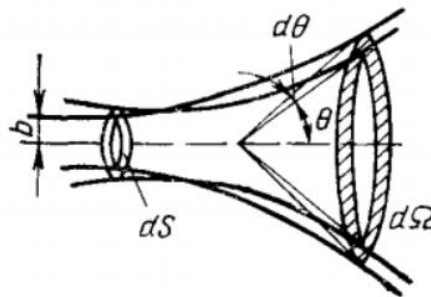


Fig. 5.5. Scattering of electrons at point defects

The scattering process is a random process, so different charge carriers deviate from their course of action at different angles. Therefore, when calculating the relaxation time, it is necessary to take into account the average value of the cross-section over the takeover angles.

Suppose the angle of scattering varies θ from $\theta + d\Omega$. Since the center of the scattering in our case has axial symmetry, these angles are the angles of two

cones, as shown in Fig. 5.5. Where $d\Omega$ denotes the spatial angle encircled between these cones

$$d\Omega = 2\pi \sin \theta d\theta \quad (5.56)$$

Sweeping processes in this case are characterized by differential shaving σ section (θ). According to (5.3), $\sigma(\theta)$ is equal to the ratio of the number of particles deviating to one centrifugal unit per second to the particle flux falling by 1 cm^2 per unit of spatial angle.

Let's assume that n electrons are moving randomly at a velocity v . If you select an area of 1 sm^2 , which is perpendicular to each direction of the electron's motion, then nv electrons fall into it in 1 second. The total number of particles deflected by an angle in 1 second within the space angle θd is equal to the flow of particles falling into the ring (Fig. 5.5). Therefore, based on the definition and (5.56) parity, we can write

$$\sigma(\theta) = \frac{nv2\pi b |db|}{nv d\Omega} = \frac{b |db|}{\sin \theta d\theta} \quad (5.57)$$

Using (5.55), the quation gets:

$$|db| = \frac{Ze^2}{2\varepsilon_r m^* v^2} \frac{1}{\sin \theta / 2} d\theta \quad (5.58)$$

We will have it. Putting (5.55) and (5.58) at (5.57),

$$\sigma(\theta) = \left(\frac{Ze^2}{2\varepsilon_r m^* v^2} \right)^2 \frac{1}{\sin^4 \theta / 2} \quad (5.59)$$

We form

(5.59) The formula was obtained by Rutherford when he was studying the scattering of alpha particles in the nuclei of the elements.

To calculate the relaxation time, it is necessary to express the probability of passing σ through the differential scattering section $\sigma(\theta)$ $W(k, k')$. Suppose that a unit volume crystal contains $N1$ input ions. Then the number of charge carriers scattered by all the centers in 1 second for each spatial angle:

$$W(k, k') nd\Omega = \sigma(\theta) N1 nv d\Omega, \quad (5.60)$$

From here

$$W(k, k') = \sigma(\theta) N1 v \quad (5.61)$$

Thus, the probability of transition from the k state to the k' state during the scattering of an electron in input ions is determined by the differential effective cross-section, the number of scattering centers, and the velocity of the charge carriers.

The collision of charge carriers with input ions can be considered elastic for two reasons. First, the mass of electrons is 104 times smaller than the mass of the ion, which is the center of scattering. Second, lattice ions are much less mobile to the equilibrium state than free charge carriers, so when a charge carrier collides with an input ion, the energy of the charge carrier changes very little.

If during an elastic collision the charge carriers are scattered into the spatial angle $d\Omega$, then the relaxation time can be written as (5.50).

$$\frac{1}{\tau(k)} = \int_{\theta} W(k, k') (1 - \cos \theta) d\Omega \quad (5.62)$$

(5.56), (5.61), (5.62),

$$\frac{1}{\tau(k)} = 2\pi N_1 v(k) \int_{\theta} \sigma(\theta) (1 - \cos \theta) \sin \theta d\theta = N_1 v(k) \sigma_c \quad (5.63)$$

It's a bo'lamiz.

$$\sigma_c = 2\pi \int_{\theta} \sigma(\theta) (1 - \cos \theta) \sin \theta d\theta \quad (5.64)$$

The mean differential effective scattering is called cross-section. The calculation of the mean on the angles θ is done by the weight function $1 - \cos \theta$. σ_c is called the effective cross-section for conductivity (or mobility). Relaxation time based on (5.63)

$$\tau(k) = \frac{1}{N_1 v(k) \sigma_c} \quad (5.65)$$

Mean free path will be equal to:

$$l(k) = \tau(k) v(k) = \frac{1}{N_1 \sigma_c} \quad (5.66)$$

It follows from (5.59) that the scattering of charge carriers by input ions is sufficiently anisotropic. Scattering dominates in small corners or large values of target distance. If we replace the value of (5.59) ($\sigma\theta$) with the formula (5.64), then the integral at the lower limit at $\theta=0$ is infinite, and the scattering section is infinite. But the scattering of charge carriers with an ionized mixture has a limited cross-section, so that the integral is performed in determining the relaxation time (5.63). The b_{max} value corresponding to θ_{min} can be determined by the following considerations: If the input ions are uniformly distributed on the crystal and their concentration is N_1 , then the average distance between them will be $N_1^{-1/3}$. The range of action of each shaver ion can be limited to half the average distance between them, so we can assume that the upper limit of the effect is $b_{max} = 1/3$. θ_{min} minimum value of min scattering angle, it can be determined from (5.55):

$$\operatorname{ctg} \frac{\theta_{\min}}{2} = \frac{\varepsilon_r m^* v^2}{2Ze^2 N_1^{1/3}} \quad (5.67)$$

$$1 - \cos \theta = 2 \sin^2 \theta / 2; \int_{\theta_{\min}}^{\pi} \frac{(1 - \cos \theta) \sin \theta d\theta}{\sin^4 \theta / 2} = -4 \sin^4 \theta_{\min} / 2 \text{ using (5.67) and}$$

(5.59) for $\sigma(\theta)$ we will have an expression.

$$\sigma_c = 2\pi \int_{\theta_{\min}}^{\pi} \sigma(\theta)(1 - \cos \theta) \sin \theta d\theta = 2\pi \left(\frac{Ze^2}{\varepsilon_r m^* v^2} \right)^2 \ln \left[1 + \left(\frac{\varepsilon_r m^* v^2}{2Ze^2 N_1^{1/3}} \right)^2 \right] \quad (5.68)$$

Based on the relaxation time of charge carriers in ionized input atoms (5.65), the equation will take on a look. This expression is called the Konuel-Weiskopf formula.

$$\tau(k) = \frac{1}{\sigma_c N_1 v(k)} = \frac{\varepsilon_r^2 m^{*2} v^3(k)}{2\pi Z^2 e^4 N_1 \ln \left[1 + \left(\frac{\varepsilon_r m^* v^2}{2Ze^2 N_1^{1/3}} \right)^2 \right]} \quad (5.69)$$

If we express velocity in terms of energy $v = (2E/m^*)^{1/2}$, the relaxation time at the scattering of the charge carriers in the input ions

$$\tau_1 = \tau_{01} E^{3/2} \quad (5.70)$$

will have the appearance, here

$$\tau_{01} = \frac{(2m^*)^{1/2} \varepsilon_r^2}{\pi Z^2 e^4 N_1 \ln \left[1 + \left(\frac{\varepsilon_r E}{Ze^2 N_1^{1/3}} \right)^2 \right]} \quad (5.71)$$

Mean free path is defined by expression:

$$l(k) = (2/m^*)^{1/2} \tau_{01} E^2 \quad (5.72)$$

Thus, the relaxation time in the dissipation of charge carriers in the input ions is proportional to the 3/2 power of energy and 1/2 of the effective mass.

5.6. Scattering in Input Atoms and Dislocations

The scattering of charge carriers by neutral mixture atoms is much weaker than scattering with input ions. However, at lower temperatures, when the concentration of atoms from ionized impurities is less than the concentration of atoms from neutral impurities, this scattering mechanism plays an important role. In this case, the scattering of the charge carrier can occur in two ways: direct elastic collision and exchange of the incoming electron with the electron of the input atom. Because of these two scattering mechanisms, determining the relaxation time is quite difficult. However, based on a hypothetical calculation that the scattering of charge carriers onto atoms of neutral compounds ε_r is the scattering of slow electrons onto a hydrogen atom located in a medium with a relative dielectric constant ε_r

$$\tau_A = \frac{e^2 m^{*2}}{20 \varepsilon_r \hbar^3 N_A} \quad (5.73)$$

The dependence is obtained, where N_A - the concentration of neutral atoms at a given temperature.

As can be seen from (5.73), the relaxation time during the scattering of charge carriers with a neutral mixture is determined by the concentration of the input atoms, not by the temperature or energy of the charge carrier being scattered. This dissolving mechanism is therefore especially important at very low temperatures, when the inlet is not ionized.

Displacement of charge carriers can occur in grid breakdowns introduced by dislocations. There is a crystal deformed sphere around the dislocation, so that the dependence of energy on the wave vector near the dislocation is different from the dispersion law in the ideal crystal. This leads to scattering of free charge carriers. In semiconductors that crystallize in a diamond-type structure, dislocations act as a chain of receptors that hold free electrons. Therefore, they represent a positive cylinder with a negative charge in the middle; that is, dislocation can be likened to a cylindrical linear negative charge with an infinitely long radius R , around which a positive spatial charge is located. Electrons moving perpendicular to dislocations cannot move freely in the direction of the field because they have to experience repulsion, bypassing the dislocations, being propagated by them. Calculations show that when charge carriers are scattered in dislocations, the relaxation time is determined not by the temperature of the crystal, but by the dislocation density pertaining to the surface of the N_D unit and the velocity v of the scattered electron according to the form expression.

$$\tau_D = \frac{3}{8 R v N_D} \quad (5.74)$$

If $v = 10^7$ cm/s; $R = 3 \cdot 10^{-5}$ cm; $N_D = 10^6$ cm⁻², the relaxation time $\tau_D = 1,25 \cdot 10^{-9}$. This value is 4-5 times greater than the relaxation time for splashing by thermal vibrations of the grille at room temperature. Therefore, at room temperature, this soaking is insignificant, but at lower temperatures, it can be noticeable.

5.7. Scattering during thermal oscillations of the grid

The interaction of an electron (or holes) with the lattice oscillations takes place in two ways: 1) the electron transfers some of its energy to the lattice, so ω_q that a certain normal oscillation of frequency q increases its quantum number by one. As a result of such a decaying process $E = \hbar \omega$, phonons with energy and quasi-impulse $Q = \hbar q$ are formed, and the number of phonons multiplies by one; Q is the

impulse that the lattice receives as a result of a collision with electrons; 2) when an electron interacts, it receives some of its energy from the lattice, so that the frequency decreases by the quantum number of a given normal oscillation. In this case, the phonon with energy and quasi-impulse disappears, and the number of phonons decreases by one. Consequently, an electron colliding with a phonon in either of these two states exchanges energy E and Q quasi-impulse with it. This scattering mechanism is called single-phonon. Since the number of phonons is determined by temperature, the scattering of electrons by the thermal oscillations of the lattice must be temperature-dependent. A multi-phonon scattering mechanism can also be found in a crystal in which several phonons are absorbed or emitted. However, the probability of such a process is much smaller compared to the probability of single-phonon scattering.

Suppose that before the collision let the electron wave vector k and energy $E(k)$, and after the collision - the wave vector k' and the energy $E(k')$. The background energy $\hbar\omega$ and the wave vector that appear or disappear during the collision have q . When an electron interacts with a phonon, it must satisfy two laws: the law of conservation of energy and the law of conservation of quasi-impulse. Therefore, when the electron absorbs the phonon

$$E'(k') = E(k) + \hbar\omega_q \quad (5.75)$$

$$k' = k + q \quad (5.76)$$

The equations become finite, the number of phonons decreases by one, i.e.

$$N'_q = N_q - 1 \quad (5.77)$$

When emitting an electronic phonon

$$E'(k') = E(k) - \hbar\omega_q \quad (5.78)$$

$$k' = k - q \quad (5.79)$$

The equations become valid and the number of phonons increases by one, i.e.

$$N'_q = N_q + 1 \quad (5.80)$$

The scattering of the charge carrier can be elastic and inelastic. To describe a splash, one needs to know the relative change of energy in a single collision

$$\delta = \Delta E / E \quad (5.81)$$

Or a scatter in a unit of time

$$A = E \delta / \tau \quad (5.82)$$

When the isotropic and elastic dissipation δ is $\ll 1$ and the mean deflection time τ is the relaxation time of the system on the impulse, $\tau E = \tau / \delta$ is the relaxation time in terms of energy.

Let us consider the scattering process of an electron in an atomic semiconductor with a cubic structure, the energy diagram of which is shown in Fig. 5.6, where a is the lattice constant in the absence of deformation. We assume that the conductivity zone of such a crystal is simple, so that the electron energy in the conductivity zone is will be equal to:

$$E = \frac{\hbar^2 k^2}{2m^*} \quad (5.83)$$

In a solid the transverse waves are shear deformation waves, and in a cubic crystal they do not cause any change in volume. The longitudinal waves, which are waves of compression and tension deformation, result in crystal size changes. As can be seen from Figure 5.6, when the crystal is compressed, the lattice constant decreases.

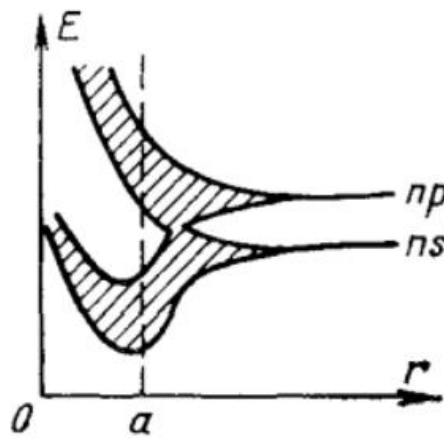


Fig. 5.6. Energy scheme of an atomically bonded semiconductor

The lower edge of the bandwidth is skewed upwards and the upper edge of the stress zone slides down, causing the bandwidth to increase. With stretching, this leads to an increase in the grid constant, while the forbidden zone decreases. Consequently, the local deformation caused by a longitudinal acoustic wave in such a crystal results in wave-like mixing of the bottom of the conduction zone and the top of the tensile zone schematically, as presented in Fig. 5.7. Therefore, the moving circuit encountered by the shear wave caused by the thermal vibrations of the lattice propagate only by longitudinal oscillations.



Fig. 5.7. The change of the bottom of the permeability zone and the ceiling of the valent zone under the influence of longitudinal oscillations of the lattice

We show that this scattering occurs in elastic and long-wave acoustic phonons. The law of conservation of energy in accordance with (5.75), (5.79) and (5.83) is written in the following form.

$$\frac{\hbar^2 k'^2}{2m^*} = \frac{\hbar^2 (k \pm q)^2}{2m^*} = \frac{\hbar^2 k^2}{2m^*} \pm \hbar \omega_q \quad (5.84)$$

$$q^2 \pm 2kq \cos \theta \mp \frac{2m^* \omega_q}{\hbar} = 0 \quad (5.85)$$

Here θ - is the angle between the vectors k and q .

In equation (5.84), the plus sign denotes the absorption state, while the minus sign denotes the phonon emission state. If an electron only interacts with long-wavelength acoustic phonons, then $\omega_q = v_t q$, where v_t is the propagation rate of the sound wave. Taking this into account, from (5.85) we find:

$$q = \mp 2k \cos \theta \pm \frac{2m^* v_t}{\hbar} \quad (5.86)$$

$$\frac{\hbar^2 k^2}{2m^*} = kT; \quad k = \sqrt{2m^* kT} / \hbar \quad \text{For being}$$

$$\frac{m^* v_t}{\hbar k} = \sqrt{\frac{T_{kr}}{T}} \quad (5.87)$$

$$T_{kr} = \frac{m^* v_t^2}{2k}$$

If we assume $v_t = 3 \cdot 10^5$ sm/s, $m^* = m_0 = 10^{-27}$ g, then $T_{kr} = 1$ K. Therefore, at temperatures much higher than 1 K (5.36), it is possible to ignore the second phenomenon. Without it

$$q = 2k \cos \theta. \quad (5.88)$$

θ Depending on the value of, the wave number can vary from $q_{min} = 0$ to $q_{max} = 2k$ that is, its average value is k . $T = 300$ K is equal to $k = k = 10^7$ cm⁻¹, and q can accept values from 0 to $2 \cdot 10^7$ cm⁻¹. This signifies a change in the energy of the phonons from 0 to $\cdot 10^{-3}$ eV. But at hona temperature, the electron energy is $\cdot 5 \cdot 10^{-2}$ eV. So that

$$\hbar \omega_q / E \approx 8 \cdot 10^{-2} \ll 1 \quad (5.89)$$

According to (3.87) and (3.96), the maximum value of the wave vector q_0 according to Debay's theory is $q_0 = 10^8$ sm⁻¹. So $q_{max} \ll q_0$. It follows that the scattering of electrons in longitudinal acoustic waves is elastic, in which case the electron absorbs or emits a phonon of $q=k$; in which case the electron is scattered by long-wave phonons and absorbs or emits a phonon of $q=k$. Since the phonon

energy is significantly less than the electron energy, $k' = k$ and $E' = E$, the phonon energy can be ignored at (5.84).

At low temperatures, the average velocity of electron motion approaches the scattering rate of the elastic waves in the crystal, and (5.84) the phonon energy cannot be ignored. In this case, a nonelastic collision of an electron with a phonon occurs.

Since the electrons are elastically scattered by longitudinal acoustic phonons, the formula (5.48) can be used to determine the relaxation time:

$$\frac{1}{\tau} = -\frac{1}{4\pi^3} \int_{V_B} W^+(k, k') \frac{q_z}{k_z} d\tau_{k'} + \frac{1}{4\pi^3} \int_{V_B} W^-(k, k') \frac{q_z}{k_z} d\tau_{k'} \quad (5.90)$$

Here the first had signifies the conquest and the second determination.

The probability of an electron *going from* state k *to state* k' when the phonon is absorbed

$$W^+(k, k') = W^+(k, k+q) = \omega(q) N_q \delta(E' - E - \hbar \nu_q) \quad (5.91)$$

expression, and when the background is emitted,

$$W^-(k, k') = W^-(k, k-q) = \omega(q) (N_q + 1) \delta(E' - E + \hbar \nu_q) \quad (5.92)$$

is defined by the expression, where $\delta(E' - E + \hbar \nu_q)$ - is the delta function representing the conservation law of energy and wave vector

$$\omega(q) = \frac{4\pi C^2 q^2}{9NM\omega_q} \quad (5.93)$$

C - is a constant with an energy dimension, N - is the number of elementary cells in a crystal; M - atomic mass.

In electron migration processes, the distribution function of phonons differs little from the balanced distribution function (3.75). So, it can be considered.

$$N_q = \frac{1}{e^{\hbar\omega_q/kT} - 1} \quad (5.94)$$

Based on (5.87) (5.94), it is possible to spread (5.94) into the row. As a result

$$\langle n \rangle = \frac{kT}{\hbar \nu_l q} \quad (5.95)$$

Taking into account (5.94) and (5.95), we put (5.91) and (5.92) on (5.90) and integrate by q

$$\tau_l = \frac{\tau_{0l}}{m^{*3/2} T E^{1/2}} \quad (5.96)$$

We'll see the results here.

$$\tau_{0l} = \frac{9\pi \hbar^4 \nu_l^2 M}{4\sqrt{2} a^3 k C^2}$$

a - lattice constant. The free path is equal to

$$l = \sqrt{2} \tau_{0l} T^{1/2} \quad (5.97)$$

Thus, in atomic semiconductors, the relaxation time in the scattering of an electron in the long-wavelength acoustic oscillations of the lattice is inversely proportional to the first degree of temperature and the energy is 1/2 of the degree, and the mean free path does not depend on the energy of the charge carriers.

Let us calculate the relaxation time of germanium in the scattering of electrons in the lattice oscillations. At room temperature, the electron mobility in pure germanium $\mu_n = 0,39 \text{ m}^2/(\text{V}\cdot\text{s})$ If we assume $m_n^* = 0,3m_0$, $e/m_n^* = 6 \cdot 10^{11} \text{ C/kg}$, then $\tau_m = m_n^* \mu_n / e = 6,5 \cdot 10^{-13} \text{ s}$.

From a comparison of the relationships (5.96) and (5.70), it follows that as the energy of the charge carriers increases, the relaxation time increases when they are scattered in an ionized mixture, and decreases when scattered in acoustic vibrations. This means that in the ingress scattering, the effect of slow electrons and spikes will be stronger, and therefore this scattering mechanism will be more pronounced at lower temperatures. When scattered in acoustic phonons, the primary contribution is made by relatively rapid electrons, and this scattering mechanism becomes more important at elevated temperatures.

In polar semiconductors, as well as in semiconductor compounds of type A^{III}B^V, where the bonding between atoms is partially ionic in nature, conductivity electrons are much more strongly affected by optical vibrations than by acoustic vibrations. Since during optical oscillations the displacement of ions in the cell occurs in the opposite direction, the opposite charges create electric fields that travel in straight waves across the crystal. The long-wavelength optical oscillations are called polarization waves. The interaction of charge carriers with polarization waves causes them to scatter. In this case, longitudinal oscillations are scattered more strongly than transverse oscillations. Each collision of an electron with an optical phonon results in the formation or loss of a phonon with energy. The frequency of the optical phonon is weakly dependent on the quasi-impulse q , so that during a collision the electron energy increases or decreases by the same amount:

$$\pm \hbar \omega_{op} = \pm \hbar \omega_0 \quad (5.98)$$

where ω_0 is the maximum frequency of the longitudinal optical network of vibration.

The energy transferred to the grid by the charge carrier in a single collision is multiplied by the ratio of the difference between the emission and absorption probability to the energy of the resulting phonon by the sum of these probabilities. But according to (5.91), the probability of absorption of the phonon is proportional to the concentration of the phonon to N_q , and according to (5.92), the probability of emission of the phonon is proportional to $N_q + 1$. This means that the energy that an electron transfers into the background is:

$$\Delta E = \hbar\omega_0 \frac{W^-(k, k') - W^+(k, k')}{W^-(k, k') + W^+(k, k')} = \hbar\omega_0 \frac{1}{2N_q + 1} = \hbar\omega_0 \frac{\hbar\omega_0}{2kT} \quad (5.99)$$

From the formula (5.99), it follows that a single collision of an electron giving energy to the phonon corresponds to a collision of $2N_q + 1$ of the phonon.

At high temperatures, when $T > \theta_0$, i.e., when the phonon energy is much smaller than the electron energy

$$\delta_{op} = \frac{(\hbar\omega_0)^2}{2(kT)^2} \ll 1 \quad (5.100)$$

becomes and can be considered elastic.

In this case, for most electrons, the energy change during absorption or emission of the optical phonon would be negligible. Therefore, we enter the relaxation time t and determine it using the formula (5.90). For the scattering of the electron in the optical phonons in the ionic crystal at higher temperatures, calculations have shown that the relaxation time is greater and a moderate free-to-run. It is determined by formulas:

$$\tau_{op} = \frac{\sqrt{2E}\hbar^2\epsilon_r}{2e^2kTm^{*1/2}} \quad (5.101)$$

$$l_{op} = \frac{\hbar^2\epsilon_r E}{e^2 m^* kT} \quad (5.102)$$

At low temperatures, in the presence of $T < \theta_0$, electrons can absorb only optical phonons, due to which their energy increases significantly. Since the probability of emitting a phonon is (5.92) greater than the probability of absorption of a phonon (5.91) $(N_q + 1)/N_q$, the time it takes for an electron to remain in a high-energy state is very short. We can hypothesize that an electron emits a phonon of the same energy immediately after absorbing a phonon. In fact, scattering can be considered as a process of phonon exchange, and so the energy of an electron does not change much, but the direction of its quasi-impulse changes strongly. This $\theta_0 = \hbar\omega/k$ makes it possible to assume that the electron interacting with the optical vibrations of the grid will be elastically scattered even at low temperatures significantly below the characteristic temperature.

At low temperatures

$$N_q = \frac{1}{e^{\hbar\omega_0/kT} - 1} = e^{-\hbar\omega_0/kT} \quad (5.103)$$

When calculating the relaxation time according to the formula (5.90), i.e., the relaxation time when only the phonon absorption is taken into account is defined by expression.

$$\tau_{op} = \frac{3\sqrt{2}\hbar^2\epsilon_r e^{\hbar\omega_0/kT}}{2e^2\sqrt{m^*\hbar\omega_0}} \quad (5.104)$$

From this you can see that the relaxation time will not depend on energy, but will depend exponentially on temperature.

In this case, the free running path of the electrons

$$l_{op} = \frac{3\hbar^2 \varepsilon_r e^{\hbar\varpi_0/kT} \sqrt{E}}{e^2 m^* \sqrt{m^* \hbar \varpi_0}} \quad (5.105)$$

It is determined by the ratio, that is, the charge is proportional to 1/2 of the energy of the carriers.

Control Questions

1. Explain the scattering mechanisms of electrons and holes?
2. Write and explain the Boltzman kinetic equation?
3. What is relaxation time?
4. Explain the scattering in the input ions?
5. What do you mean by scattering in dislocations?
6. Explain scattering in thermal oscillations of the grid.
7. What is a shaving section?

Chapter VI

KINETIC PHENOMENA IN SEMICONDUCTORS

6.1. Unbalanced Distribution Function

To describe escape phenomena using Boltzman's kinetic equation, one needs to know the unbalanced distribution function of electrons and canines. We define it for an unturned semiconductor having both electric and magnetic fields and a temperature gradient. Energy of electrons for spherical isoenergetic surface semiconductor

$$E = \frac{\hbar^2 k^2}{2m_n} \quad (6.1)$$

Their speed

$$v = \frac{e}{\hbar} \frac{\partial E}{\partial \mathbf{k}} = \frac{\hbar}{m_n^*} \mathbf{k} \quad (6.2)$$

Electron in the presence of electric and magnetic fields

$$\mathbf{F} = -e(\mathcal{E} + [\mathbf{v}\mathbf{B}]) \quad (6.3)$$

(5.44) kinetic equation for force effects and steady-state

$$(\mathbf{v}\nabla_{\mathbf{r}}f) - \frac{e}{\hbar}(\boldsymbol{\delta} + [\mathbf{v}\mathbf{B}])\nabla_{\mathbf{k}}f = -f_1(\mathbf{k})/\tau_e \quad (6.4)$$

takes the view, where τ_e – the time of relaxation of the electron; φ – electrostatic potential. $\tau_e E = -\nabla\varphi$; φ

We assume that the non-negotiable distribution function is less different from the balanced distribution function. Then in equation (6.4), in all cases where the magnetic field is not involved, $f(\mathbf{k}) = f_0 + f_1(\mathbf{k})$ f can be replaced by f_0 .

$$(\mathbf{v}\nabla_{\mathbf{r}}f_0) - \frac{e}{\hbar}(3\nabla_{\mathbf{k}}f_0) - \frac{e}{\hbar}[\mathbf{v}\mathbf{B}]\nabla_{\mathbf{k}}f = -\frac{f_1(\mathbf{k})}{\tau_e} \quad (6.5)$$

For the unsaturated electron gas

$$\frac{\partial f_0}{\partial E} = -\frac{1}{kT} e^{-(E-F)/kT} \quad (6.6)$$

$$\frac{\partial f_0}{\partial T} = \frac{E-F}{kT^2} e^{-(E-F)/kT} = \frac{F-E}{T} \frac{\partial f_0}{\partial E} \quad (6.7)$$

$$\frac{\partial f_0}{\partial F} = \frac{1}{kT} e^{-(E-F)/kT} = -\frac{\partial f_0}{\partial E} \quad (6.8)$$

So that

$$\nabla_{\mathbf{r}}f_0 = \frac{\partial f_0}{\partial T}\nabla T + \frac{\partial f_0}{\partial F}\nabla F = \frac{\partial f_0}{\partial E}\left(\frac{F-E}{T}\nabla T - \nabla F\right) \quad (6.9)$$

$$\nabla_{\mathbf{k}}f_0 = \frac{\partial f_0}{\partial E} \frac{\hbar^2}{m_{rt}^*} \mathbf{k} = \hbar \frac{\partial f_0}{\partial E} \mathbf{v} \quad (6.10)$$

and

$$[\mathbf{v}\mathbf{B}]\nabla_{\mathbf{k}}f = [\mathbf{v}\mathbf{B}]\nabla_{\mathbf{k}}(f_0 + f_1) = \hbar \frac{\partial f_0}{\partial E} [\mathbf{v}\mathbf{B}]\mathbf{v} + [\mathbf{v}\mathbf{B}]\nabla_{\mathbf{k}}f_1 \quad (6.11)$$

But that's why $[\mathbf{v}\mathbf{B}]\mathbf{v} = [\mathbf{v}\mathbf{v}]\mathbf{B} = 0$

$$f_1(\mathbf{k}) = -\frac{\partial f_0}{\partial E} \chi_{vi} \mathbf{k} = -\frac{m_{nt}^*}{\hbar} \frac{\partial f_0}{\partial E} (\chi_{rs} \mathbf{v}) \quad (6.12)$$

$$\nabla_{\mathbf{k}} f_1 = -\nabla_{\mathbf{k}} \left\{ \frac{\partial f_0}{\partial E} (\chi_n \mathbf{k}) \right\} = -\frac{\partial f_0}{\partial E} \chi_n - \left\{ \hbar \mathbf{k} \frac{\partial}{\partial E} \left(\frac{\partial f_0}{\partial E} \chi_n \right) \right\} \mathbf{v} \quad (6.13)$$

Without it

$$(\mathbf{v}\mathbf{B})\nabla_{\mathbf{k}} f = [\mathbf{v}\mathbf{B}]\nabla_{\mathbf{k}} f_1 = -\frac{\partial f_0}{\partial E} ([\mathbf{v}\mathbf{B}]\chi_n) = -\frac{\partial f_0}{\partial E} [\mathbf{B}\chi_n] \mathbf{v} \quad (6.14)$$

Because $[\mathbf{v}\mathbf{B}]\chi_n = [\mathbf{B}\chi_n]\mathbf{v}$

using (6.8), (6.9), (6.10) and (6.14) and (6.5)

$$\begin{aligned} f_1(\mathbf{k}) &= -\frac{m_n^*}{\hbar} \frac{\partial f_0}{\partial E} (\chi_n \mathbf{v}) = \\ &= -\tau_e \frac{\partial f_0}{\partial E} \left\{ \frac{F-E}{T} \nabla T - \nabla(F - e\varphi) + \frac{e}{\hbar} [\mathbf{B}\chi_n] \right\} \mathbf{v} \end{aligned} \quad (6.15)$$

We take it. Based on (6.12), respectively

$$\chi_n = -\frac{\hbar}{m_n^*} \tau_e \left\{ \frac{E-F}{T} \nabla T + \nabla(F - e\varphi) - \frac{e}{\hbar} [\mathbf{B}\chi_n] \right\} \quad (6.16)$$

E' energiya va Fermi sathi F'

$$\begin{aligned} E' &= -\frac{\hbar^2 \mathbf{k}'^2}{2m_p^*} \\ F' &= -E_g - F \end{aligned}$$

is equal to the width of the prohibited zone, E_g

$$\chi_P = -\frac{\hbar}{m_p^*} \tau_h \left\{ \frac{E'+E_g+F}{T} \nabla T - \nabla(F - e\varphi) + \frac{e}{\hbar} [\mathbf{B}\chi_P] \right\}. \quad (6.19)$$

It is the time of relaxation of the hives. τ_h

Identification of (6.16) and (6.19) $\chi_n \chi_p$

$$\mathbf{x} = \mathbf{a} + [\mathbf{b}\mathbf{x}]. \quad (6.20)$$

The vector is brought to the solution of the equation. (6.20) We multiply the equation by scalar to the vector: $\mathbf{b}$

$$(bx) = (ba) + b[bx] = (ba) \quad (6.21)$$

Now (6.20) we multiply the equation vector by vector: $\mathbf{b}$

$$[\mathbf{b}\mathbf{x}] = [\mathbf{b}\mathbf{a}] + [\mathbf{b}[\mathbf{b}\mathbf{x}]] = [\mathbf{b}\mathbf{a}] + \mathbf{b}(\mathbf{b}\mathbf{x}) - \mathbf{x}(\mathbf{b}\mathbf{b}) = [\mathbf{b}\mathbf{a}] + \mathbf{b}(\mathbf{b}\mathbf{a}) - \mathbf{x}b^2 \quad (6.22)$$

$$\mathbf{x} = \mathbf{a} + [\mathbf{b}\mathbf{a}] + \mathbf{b}(\mathbf{b}\mathbf{a}) - \mathbf{x}b^2 \quad (6.23)$$

From here

$$\mathbf{x} = \frac{\mathbf{a} + [\mathbf{b}\mathbf{a}] + \mathbf{b}(\mathbf{b}\mathbf{a})}{1+b^2} \quad (6.24)$$

If we use (6.24), (6.16) and (6.19) can be written as follows

$$\chi_n = -\frac{\hbar\tau_e}{m_n^*} \left[\frac{\frac{E-F}{T}\nabla T + \nabla(F-e\varphi) + \frac{e\tau_e}{m_n^*} \left[\mathbf{B} \left\{ \frac{E-F}{T}\nabla T + \nabla(F-e\varphi) \right\} \right]}{1 + \left(\frac{e\tau_e}{m_n^*} B \right)^2} + \right. \\ \left. + \frac{\left(\frac{e\tau_e}{m_n^*} \right)^2 \left(\mathbf{B} \left\{ \frac{E-F}{T}\nabla T + \nabla(E-e\varphi) \right\} \right) \mathbf{B}}{1 + \left(\frac{e\tau_e}{m_n^*} B \right)^2} \right]; \quad (6.25)$$

$$\chi_p = -\frac{\hbar\tau_h}{m_p^*} \left[\frac{\frac{E'+Eg+F}{T}\nabla T - \nabla(F-e\varphi) - \frac{e\tau_h}{m_p^*} \left[\mathbf{B} \left\{ \frac{E'+Eg+F}{T}\nabla T - \nabla(F-e\varphi) \right\} \right]}{1 + \left(\frac{e\tau_h}{m_p^*} B \right)^2} + \right. \\ \left. + \frac{\left(\frac{e\tau_h}{m_p^*} \right)^2 \left(\mathbf{B} \left\{ \frac{E'+Eg+F}{T}\nabla T - \nabla(F-e\varphi) \right\} \right) \mathbf{B}}{1 + \left(\frac{e\tau_h}{m_p^*} B \right)^2} \right]; \quad (6.26)$$

6.2. Specific Electrical Conductivity of Semiconductors

In the absence of a temperature gradient and magnetic field, we calculate the specific electrical conductivity of a homogeneous unseparated semiconductor whose isoenergetic surface is a sphere. Let the semiconductor be located in a homogeneous electric field,

Number of electrons per unit size crystal based on volume element (5.12)

$d\tau_{\mathbf{k}}$

$$dn = f(\mathbf{k}) \frac{d\tau_{\mathbf{k}}}{4\pi^3} \quad (6.27)$$

This is when electrons travel with velocities $\mathbf{v}$ under the action of an external electric field

$$d\mathbf{J} = -e\mathbf{v}f(\mathbf{k}) \frac{d\tau_{\mathbf{k}}}{4\pi^3} \quad (6.28)$$

It forms an elemental current density. If a semiconductor has two different charge carriers, the total current density is

$$\mathbf{J} = \mathbf{J}_n + \mathbf{J}_p = -\frac{e}{4\pi^3} \int_{(V_E)} \mathbf{v}_n f_n(\mathbf{k}) d\tau_{\mathbf{k}} + \frac{e}{4\pi^3} \int_{(V_E)} \mathbf{v}_p f_p(\mathbf{k}') d\tau_{\mathbf{k}'}, \quad (6.29)$$

where $\mathbf{v}_n$ and $\mathbf{v}_p$ - velocities and non-equilibrium distribution functions of electrons and pinnies. Integration should be carried out over the entire Brilllueen zone. $\mathbf{v}_n, f_n(\mathbf{k})$ va $\mathbf{v}_p, f_p(\mathbf{k}')$

If we put the balanced distribution functions in (6.29),

$$\mathbf{J} = \int_{(V_E)} \mathbf{v} f_0(\mathbf{k}) d\tau_{\mathbf{k}} = 0 \quad (6.30)$$

It turns out that where $f_0(\mathbf{k})$ is a double function, $f_0(\mathbf{k})$ is zero when integrated with respect to the odd function and symmetric boundaries. $\mathbf{v} f_0(\mathbf{k}) = \frac{\hbar}{m^*} \mathbf{k} f_0(\mathbf{k})$

$$f_n(\mathbf{k}) = f_0 + f_1^{(n)}(\mathbf{k}) = f_0 - \frac{\partial f_0}{\partial E}(\chi_n \mathbf{k}) \quad (6.31)$$

Considering (6.30), instead of (6.29)

$$\begin{aligned} \mathbf{J} &= -\frac{e}{4\pi^3} \int_{(V_B)} \mathbf{v}_n f_1^{(n)}(\mathbf{k}) d\tau_{\mathbf{k}} + \frac{e}{4\pi^3} \int_{(V_B)} \mathbf{v}_p f^{(p)}(\mathbf{k}') d\tau_{\mathbf{k}'} = \\ &= \frac{e\hbar}{4\pi^3 m_n^*} \int_{(V_B)} \frac{\partial f_0}{\partial E}(\chi_n \mathbf{k}) \mathbf{k} d\tau_{\mathbf{k}} - \frac{e\hbar}{4\pi^3 m_p^*} \int_{(V_B)} \frac{\partial f_{0p}}{\partial E'}(\chi_p \mathbf{k}') \mathbf{k}' d\tau_{\mathbf{k}'}. \end{aligned} \quad (6.32)$$

We will have it. Since the semiconductor is homogeneous and the temperature gradient is absent, ($\nabla T = 0$) (6.9) is based on the equation $\nabla_r f = 0 \nabla T = 0$

$$\nabla F = 0 \quad (6.33)$$

In the presence of an electric field, $E = -\nabla\varphi$ and in the absence of a magnetic field ($\mathbf{B}0$), (6.33), from (6.25) and (6.26)

$$-\chi_n = -\frac{e\hbar\tau_e}{m_n^*} E \quad (6.34)$$

$$\chi_p = \frac{e\hbar\tau_h}{m_p^*} \quad (6.35)$$

will have expressions. Using (6.34), for the circuit current density, we write the following

$$\mathbf{J}_n = -\frac{e^2 \hbar^2}{4\pi^3 m_n^{*2}} \int_{(V_E)} \tau_e \frac{\partial f_0}{\partial E}(\mathbf{k}E) \mathbf{k} d\tau_{\mathbf{k}}. \quad (6.36)$$

In the calculation from the bottom of the energy conductivity field for an unfused semiconductor, the number of electrons based on (4.66)

$$n = N_c e^{-(E_c - F)/kT} = N_c e^{F/kT} \quad (6.37)$$

It is determined by attitude. For being for an undisturbed state $f_0 = e^{-(E-F)/kT}$

$$\frac{\partial f_0}{\partial E} = -\frac{1}{kT} e^{-(E-F)/kT} = -\frac{n}{kT N_c} e^{-E/kT}. \quad (6.38)$$

Equation (6.36) is applied on (6.38):

$$\mathbf{J}_n = \frac{e^2 \hbar^2 n}{4\pi^8 m_n^{*2} kT N_c} \int_{(V_E)} \tau_e e^{-E/kT} (\mathbf{k}\boldsymbol{\varepsilon}) \mathbf{k} d\tau_{\mathbf{k}} \quad (6.39)$$

$$\int F(\mathbf{k})(\mathbf{a}\mathbf{k}) \mathbf{k} d\tau_{\mathbf{k}} = \frac{4\pi}{3} \mathbf{a} \int_0^\infty F(\mathbf{k}) \mathbf{k}^4 d\mathbf{k} \quad (6.40)$$

If we use the attitude, then the (6.39) can write in the form. , defining that, $\alpha = E/kT$

$$\mathbf{J}_n = \frac{e^2 \hbar^2 n \boldsymbol{\varepsilon}}{4\pi^3 m_n^* kT N_c} \frac{4\pi}{3} \int_0^\infty \tau_e e^{-E/kT} \mathbf{k}^4 d\mathbf{k} \quad (6.41)$$

$$\mathbf{k}^4 d\mathbf{k} = \frac{(2m_n^* kT)^{5/2}}{2\hbar^5} \alpha^{3/2} d\alpha \quad (6.42)$$

We will have it and taking into account that (6.42)(6.41) $N_c = 2(2\pi m_n^* kT/\hbar^2)^{3/2}$

$$\mathbf{J}_n = \frac{e^2 n}{m_n^*}, \boldsymbol{\varepsilon} \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e e^{-\alpha} \alpha^{3/2} d\alpha \quad (6.43)$$

We enter a mark

$$\langle \tau_e \rangle = \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e e^{-\alpha} \alpha^{3/2} d\alpha \quad (6.44)$$

Without it

$$J_n = \frac{e^2 n \langle \tau_e \rangle}{m_n^*} \mathcal{B} = en\mu_n \mathcal{E} = \sigma_n \mathcal{B} \quad (6.45)$$

here

$$\mu_n = \frac{e \langle \tau_e \rangle}{m_n^*} \quad (6.46)$$

-It is obtained for its activeness,

$$\sigma_n = en\mu_n \quad (6.47)$$

- Specific conductivity for electrons.

Spurs of the Spurs, respectively

$$J_p = \frac{e^2 p \langle \tau_h \rangle}{m_p^*} \mathcal{E} = ep\mu_p \mathcal{E} = \sigma_p \mathcal{E} \quad (6.48)$$

Here is the average hives rest time

$$\langle \tau_h \rangle = \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_h e^{-\alpha} \alpha^{3/2} d\alpha \quad (6.49)$$

Drift mobility of hives

$$\mu_p = \frac{e \langle \tau_h \rangle}{m_p^*} \quad (6.50)$$

and specific conductivity

$$\sigma_p = ep\mu_p \quad (6.51)$$

The full current density in a semiconductor with two different charge carriers, respectively

$$\mathbf{J} = \mathbf{J}_n + \mathbf{J}_p = (\sigma_n + \sigma_p) \mathbf{E} = \sigma \mathbf{E} \quad (6.52)$$

This is the specific conductivity of a semiconductor. $\sigma = \sigma_n + \sigma_p$ –

6.3. The dependence of the mobility of charge carriers on temperature

It follows from (6.46) and (6.50) that it is necessary to know the average relaxation time to calculate the mobility of charge carriers. But the time of relaxation is strongly influenced by the digestion mechanisms. Therefore, in the collision of charge carriers with various defects in the Crystal Grid, the mobility is invariably dependent on the temperature.

A. Scattering on grid oscillations

Relaxation time based on (5.96) in scattering of charge carriers in lattice acoustic oscillations in atomic semiconductors will be equal to (6.44)

$$\tau_l = \frac{\tau_0 t}{m^{*3/2}} T^{-1} E^{-1/2} \quad (6.53)$$

$$\langle \tau_l \rangle = \frac{4}{3\sqrt{\pi}} \frac{\tau_0 t}{m^{3/2} k^{1/2}} T^{-3/2} \quad (6.54)$$

and

$$\mu_l = \frac{4e}{3\sqrt{\pi} \pi} \frac{\tau_0 t}{k^{1/2} m^{3/2}} T^{-3/2} \quad (6.55)$$

Mobility can be expressed by the free-treadmill. Using (5.97), and we have following equation:

$$\mu_l = \frac{4el}{3(2\pi m^* kT)^{1/2}} \quad (6.56)$$

From (6.55) it follows that in atomic semiconductors, the mobility of charge carriers decreases with temperature increase when scattered in the thermal oscillations of the lattice. In addition, mobility is inversely proportional to 5/2 of the effective mass. Since it is, the mobility of electrons is greater than that of the canines. In small effective masses of charge carriers, mobility is at its highest. (Table 6.1). $m_n^* < m_p^*$

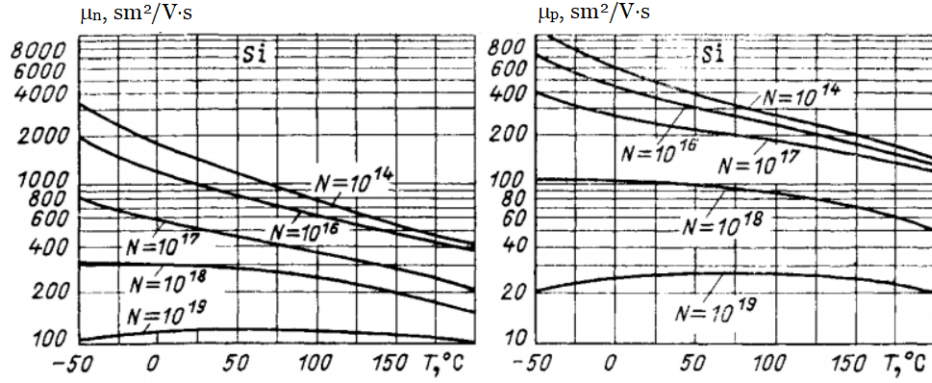


Fig. 6.1. Temperature dependence of electrons and holes drift mobility in silicon

B. Scattering in ionized inlets. Based on the relaxation rate (5.70) and (5.71)

$$\tau_I = \frac{\sqrt{2}\epsilon_r^2 m^{*1/2} E^{3/2}}{\pi Z^2 e^4 N_I \ln \left[1 + \left(\epsilon_r E / Ze^2 N_I^{1/3} \right)^2 \right]} \quad (6.57)$$

In the averaging of this expression, it can be deduced from the integral, since the logarithmic has change slowly when the energy that reaches the maximum is released. It has to be for that. Therefore, the average relaxation time $E = 3kT$

$$\langle \tau_I \rangle = \frac{8\sqrt{2}\epsilon_r^2 k^{3/2} m^{*1/2} T^{3/2}}{\pi^{3/2} Z^2 e^4 N_I \ln \left[1 + \left(3\epsilon_r kT / Ze^2 N_I^{1/3} \right)^2 \right]} \quad (6.58)$$

Movement is:

$$\mu_I = \frac{8\sqrt{2}\epsilon_r^2 k^{3/2} T^{3/2}}{\pi^{3/2} Z^2 e^3 N_I m^{*1/2} \ln \left[1 + \left(3\epsilon_r kT / Ze^2 N_I^{1/3} \right)^2 \right]} \quad (6.59)$$

At extremely high temperatures can be considered as:

$$\mu_I \approx \mu_{0I} T^{3/2} \quad (6.60)$$

From the formula (6.60), it follows that the mobility of the charge carriers resulting from scattering into the input ions decreases with the decrease in temperature. This is because with increasing crystal temperature, the heat rate of the charge carriers increases, and this leads to a weakening of the interaction of the charge carriers with the atoms of ionized mixtures, since the duration of the interaction is reduced. In addition, with an increase in the inlet concentration at a

certain temperature, mobility decreases, which increases the effect of decubation. This is consistent with experimental data. As an example, Fig. 6.1 shows the temperature dependence of electron and pit mobility in silicon samples with different concentrations of different compounds.

C. Scattering in neutral inputs

When the charge carriers are scattered at neutral inputs, according to the expression (5.73), the relaxation time does not depend on the temperature and energy of the charge carrier. Therefore, the mobility of the charge carriers, caused by the dispersion of neutral substances in atoms, does not depend strictly on temperature

$$\mu_A = e^2 m^* / 20 \varepsilon_r \hbar^3 N_A. \quad (6.61)$$

There is no ballpark.

G. Scattering in dislocations

Mobility when charge carriers are scattered in dislocations will be equal to:

$$\mu_D = \frac{e T^{-1/2}}{2^{3/2} \pi^{1/2} k^{1/2} m^{*1/2} R N_D} \quad (6.62)$$

D. Complex Digestion Mechanism

In true semiconductors, the constant circuiticity of the lattice field is broken by the presence of various defects in it. Therefore, the seething mechanism is complicated. If we consider all the scattering mechanisms to be independent, then the total probability of scattering is equal to the sum of the probability of scattering at the center of each scattering, i.e.,

$$W(k, k') = \sum_i W_i(k, k') \quad (6.63)$$

In this case, the total relaxation time is based on (5.49) τ

$$\tau^{-1} = \frac{1}{4\pi^3} \int_{(v_5)} \sum_i W_i(\mathbf{k}, k') (1 - \cos \theta) d\tau_{k'} = \sum_i^1 \tau_i^{-1} \quad (6.64)$$

or

$$\tau = (\sum_i \tau_i^{-1})^{-1} \quad (6.65)$$

Naturally, the contribution of various settling mechanisms to the overall relaxation time depends on the crystal temperature. When the energy of the charge carriers changes, the role of some scattering mechanisms decreases, while the role of others increases, so the nature of the motion of the charge carriers under the influence of an external electric field depends on which of the scattering mechanisms is dominant over a given temperature range. At very low temperatures (6.65), expression can only be accounted for scattering by input atoms and dislocations. With increasing temperature, the role of these mechanisms decreases relative to their absorption by ingress ions. At high temperatures, the scattering of phonons is dominant. Thus, the total mobility is determined by scattering into input ions and

phonons, and according to (6.65) μ it can be assumed that $\mu = (\mu_1^{-1} + \mu_2^{-1})^{-1}$, μ and depending on the temperature, given (6.55) and (6.59), the form of that relationship is determined. It is defined as:

$$\frac{1}{\mu} = aT^{-3/2} + bT^{3/2} \quad (6.66)$$

where a and b are constants.

(6.66) follows from the formula that with the increase in temperature in atomic semiconductors, the mobility of the charge carriers increases proportionally to $T^{3/2}$, if scattering occurs only in mixture ions, it will pass the maximum, and if the scattering centers are only long p-waves, it will decrease proportionally to $T^{3/2}$.

The experiment shows that such a dependence of mobility on temperature is indeed observed in semiconductors. As an example, Fig. 6.2 shows the temperature-dependent relationship of electron and hole mobility in the indium antimonide.

E. Dependence of Specific Conductivity on Temperature

The specific conductivity of a private semiconductor can be determined by considering that the mobility and concentration of charge carriers depend on temperature. The equation can write in the following form:

$$\sigma = en_i\mu_n + ep_i\mu_p = AT^{3/2+p}e^{-E_i/2kT} \quad (6.67)$$

Multiplication changes slowly with temperature changes, while multiplication is strongly dependent on temperature. Therefore, at insufficiently high temperatures $E_\Omega \gg kT$ and $AT^{3/2+p}e^{-E_i/2kT}$

$$AT^{3/2+p} = \sigma_0 = \text{const},$$

and (6.67) the expression can be replaced by a simpler expression.

$$\sigma = \sigma_0 e^{-E_\ell/2kT} \quad (6.69)$$

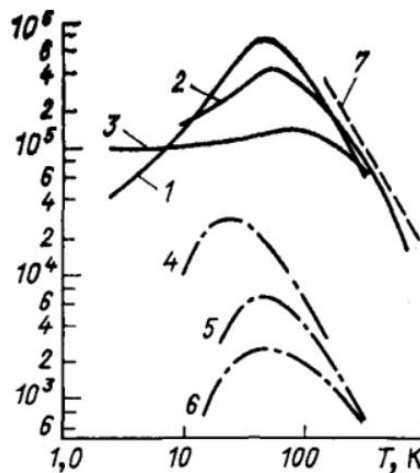


Fig. 6.2. Temperature dependence of electrons and canine mobility in indium antimonide

As can be seen from Fig. 6.3, the semiconductor conductivity of the semiconductor at sufficiently high temperatures in the input semiconductor will be specific. At low temperatures, the concentration of the main charge carrier in a

poorly spaced area is constant, and the electrical conductivity changes with the change in temperature depending on the mobility.

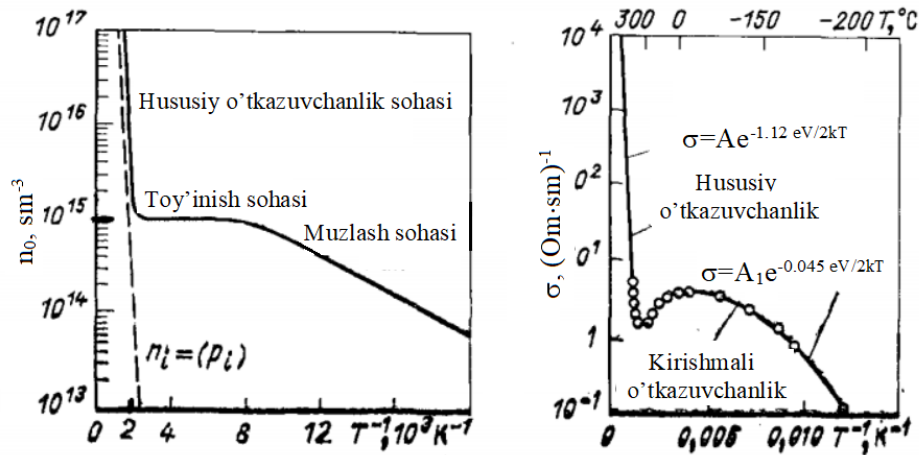


Fig 6.3. Dependence of charge carrier concentration and conductivity on the harat

6.4. Hall effect

Kinetic phenomena when a conductor is simultaneously exposed to electric and magnetic fields are called galvanomagnetic phenomena. The Hall effect is one such phenomenon. For the weak magnetic field, consider the Hall effect. Weak magnetic field is an area in which the cycle of rotation in the magnetic field is much greater than the relaxation time.

$$\tau \ll T_c \quad (6.70)$$

m^* The frequency of rotation of the effective mass charge carrier, i.e., the cyclotron frequency resonance (2.150), is based on:

$$\omega_c = \frac{2\pi}{T_c} = \frac{e}{m^*} B \quad (6.71)$$

is equal to . Therefore, in weak areas

$$\frac{\tau \omega_c}{2\pi} = \frac{e\tau}{2\pi m^*} B = \frac{\mu B}{2\pi} \ll 1, \quad (6.72)$$

And in strong areas

$$\frac{\tau \omega_c}{2\pi} = \frac{\mu B}{2\pi} \gg 1. \quad (6.73)$$

Thus, the condition of strong and weak fields is determined not only by magnetic field induction, but also by mobility.

If we use the free treadmill and the radius of the cyclotron orbit , $l = v\tau = v/\omega_c$

$$l/r = \tau \omega_c. \quad (6.74)$$

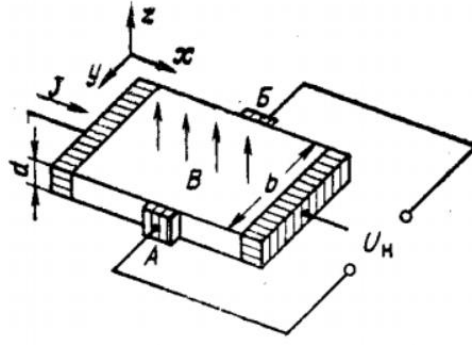


Fig. 6.4. Hall effect measurement sample

Corininth says that in a weak magnetic field, $B r \gg l$ travels a short distance until the charge carrier collides while moving in a circle in a steep plane to B, while in a strong magnetic field it travels a great distance and bends violently.

Suppose that a dam is under the action of an electric field from a separate conductor in the form of a rectangular plate

$$\mathbf{J} = -en\mathbf{v} = \sigma\mathbf{E}. \quad (6.75)$$

vine density is flowing. If the semiconductor is homogeneous, the equivalent surfaces are perpendicular to the direction of the electric field and, respectively, to the current density vector. Therefore, the difference in potentials between the points in Figure 6.4 is zero.

We now place the semiconductor on a magnetic field perpendicular to the current vector, as shown in Figure 6.4. In this case, the charge carrier moving at velocity $\mathbf{v}$ is affected by the Lorentz force

$$\mathbf{F} = \pm e[\mathbf{v}\mathbf{B}], \quad (6.76)$$

$\mathbf{v}$ and $\mathbf{B}$ are the force and perpendicular where plus refers to the chapel, minus – to the electron.

But

$$\mathbf{v} = \pm \frac{e\langle\tau\rangle}{m^*} \mathbf{E}, \quad (6.77)$$

So that

$$\mathbf{F} = \frac{e^2\langle\tau\rangle}{m^*} [\mathbf{E}\mathbf{B}] \quad (6.78)$$

It follows from (6.78) that the direction of the Lorentz force is determined not by the sign of the charge carrier, but by the direction of the vectors $\mathbf{E}$ and $\mathbf{B}$ or $\mathbf{J}$ and $\mathbf{B}$. Consequently, if the velocity of the charge carriers is determined by the external electric field, then under the action of the Lorentz force, the electrons and the fins spin in the same direction.

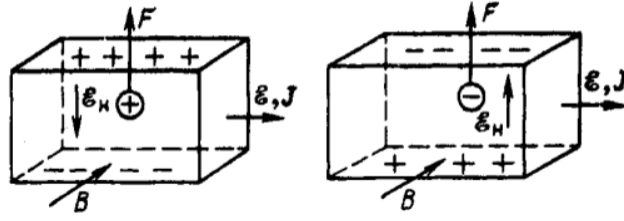


Fig. 6.5. Deflection of the magnetic field of the current carriers in p-type and n-type samples.

For the selected B and E directions shown in Fig. 6.5, the Lorentz force is directed upwards F . Under the influence of this force, the holes in the acceptor semiconductor (Fig. 6.5, a) and the electrons in the donor semiconductor (Fig. 6.5, b) are pushed onto the upper surface of the sample, resulting in their deficiency on the lower surface of the sample, resulting in an opposite pointed charge on the upper surface. By the separation of charge occurs an electric field of intensity E_H that is perpendicular to the direction of the magnetic field. The direction of this field, called Hall Square, depends on the mark of the chargers. In our case, the E_H Hall field is directed downwards in the p-sample and upwards in the n-sample. The phenomenon of the appearance of a transverse electric field in a semiconductor where a current passes through under the influence of a magnetic field is called the Hall effect. The field strength E_H increases until the force arising from this field compensates for the Lorentz force:

$$-eE_x = evB \quad (6.79)$$

In this case, the charge carriers move along the sample only longitudinally under the influence of the electric field E , and therefore the current density in the direction corresponds to the force J .

F. The total electric field voltage vector $E' = E + E_H$ is inverted by a certain angle, as shown in Fig. 6.6. J and E are called Hall's angle with respect to the direction of the angle's flow φ . The equipotential surfaces are now φ rotated at an angle relative to the initial position, so that a potential difference between points A and B, called the Hall EYUK, appears in the sample (Figure 6.4). If the width of the prayer is b ,

$$V_H - \therefore - \quad (6.80)$$

Using (6.75), we will have following:

$$V_H = -\frac{1}{en}JBb = RJBb. \quad (6.81)$$

The magnitude R in (6.81) is called the Hall constant. Hall's constant in the case of electrons is equal to .

$$R = -\frac{1}{en^2} \quad (6:82)$$

If there are holes with a concentration of charge carriers p , e , so $e\mathcal{E}_H = evB$

$$R = \frac{1}{ep} \quad (6.83)$$

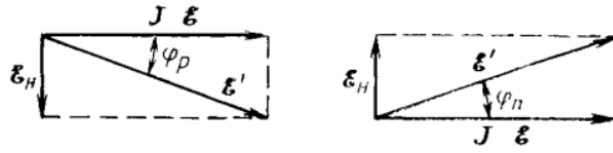


Fig. 6.6. Hall angle in p-type and n-type samples.

Based on equations (6.82) and (6.83), the Hall coefficient is inversely proportional to the concentration of charge carriers, and its sign is consistent with the charge carrier sign. When considering the Hall effect, the statistical distribution of charge carriers by energy and the dependence of relaxation time on energy were not taken into account.

Hall angle (Fig. 6.6)

$$\tan \varphi = \frac{B_H}{\mathcal{E}} \quad (6.84)$$

can be determined from the attitude, the sign of the Hall field is determined by the direction: positive for the holes, negative for the electrons. If the magnetic field is weak, $\tan \varphi$

$$\varphi = \frac{\mathcal{F}_H}{\mathcal{E}} = \frac{RJB}{\mathcal{E}} = R\sigma B, \quad (6.85)$$

It can be considered. From here

$$R = \frac{\varphi}{\sigma B}. \quad (6.86)$$

It turns out that there is a lot of work to be done. For Electronic Semiconductor

$$\sigma_n = en\mu_n, \quad (6.87)$$

Hall's Angle for Division

$$\varphi_n = R_n \sigma_n B = \mu_n B. \quad (6.88)$$

Same for hole semiconductor

$$\varphi_p = R_p \sigma_p B = \mu_p B \quad (6.89)$$

6.5. Hall effect in semiconductors of two different types

Suppose that the current concentration and mobility in a semiconductor n, μ_n and p, μ_p carried out by electrons and holes that are equal to and equal to the current in a semiconductor, respectively. We determine the Hall coefficient for the weak magnetic field state. Suppose that the electric field voltage vector E , the current

density vector J , and the magnetic induction vector B are oriented with respect to the coordinate axes, as shown in Figure 6.7.

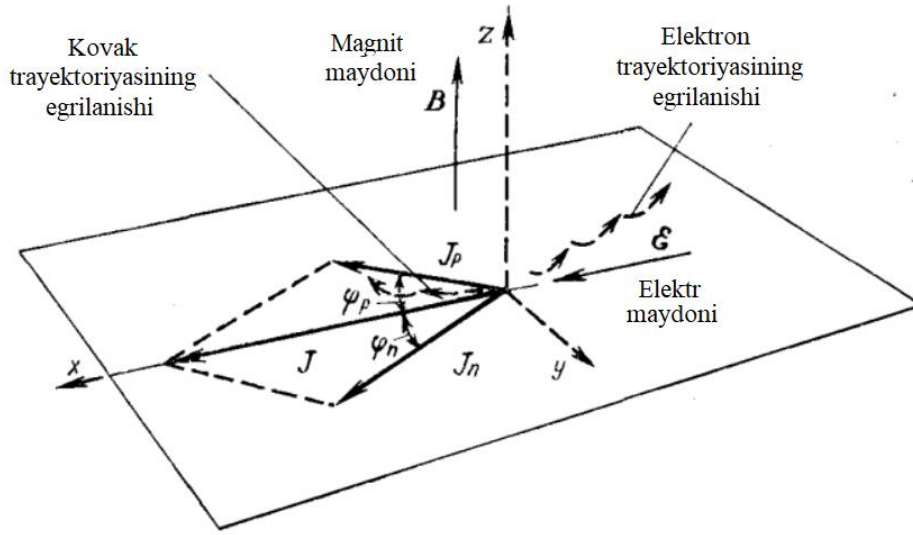


Fig. 6.7. The Hall effect in the presence of two different types of charge carriers

Under the action of electric field, charge carriers accelerate. As the drift's velocity increases, the charge carrier is affected by the Lorentz turning force. When exposed to a magnetic field, the trajectory of the charge carrier is curved until it collides with defects in the crystal lattice (the magnetic field is weak and $l \ll r$). During a collision, the charge carrier transfers the energy accumulated in the electric field to the grid and then acts again under the action of the E field. As a result, the density of the pit current deviates by an angle from the direction of the electric field J_p . Under the influence of a magnetic field, electrons are deflected in the same direction as the spurs, and the electron current density φ_p is J_n deflected from the direction of the electric field to an angle in the opposite direction (Fig. 6.7). In this case, if we use the relation (6.32), the total current density is equal to: φ_n

$$\mathbf{J} = \frac{e\hbar}{4\pi^3 m_n^*} \int_{(V_E)} \frac{\partial f_0}{\partial E} (\chi_n \mathbf{k}) \mathbf{k} d\tau_{\mathbf{k}} - \frac{e\hbar}{4\pi^3 m_p^*} \int_{(V_E)} \frac{\partial f_{0p}}{\partial E'} (\chi_p \mathbf{k}') \mathbf{k}' d\tau_{\mathbf{k}'} \quad (6.90)$$

We calculate the electron and cavity current densities for a homogeneous semiconductor having spherical isoenergetic surfaces located in a weak magnetic field where the temperature gradient is absent. and because of the fact that (6.25) and (6.26) For weak magnetic field it comes to the view:

$$\nabla_r f = \frac{E-F}{T} \nabla T + \nabla F = 0, \quad \mathcal{E} = -\nabla \gamma \frac{e\tau}{m^*} \dot{B} = \mu B \ll 1$$

$$\chi_n = -\frac{e\hbar\tau_e}{m_n^*} \left\{ \mathbf{B} + \frac{e\tau_e}{m_n^*} [\mathbf{B}\mathbf{B}] \right\} \quad (6.91)$$

$$\chi_D = \frac{e\hbar\tau_h}{m_p^*} \left\{ \mathbf{B} - \frac{e\tau_h}{m_p^*} [\mathbf{B}\mathbf{B}] \right\} \quad (6.92)$$

Now, using (6.91), we find the circuit collector of the current:

$$\mathbf{J}_n = -\frac{e^2 \hbar^3}{4\pi^3 m_n^*} \int_{(V_B)} \frac{[\partial f_0]}{\partial E} \left\{ \tau_c(\mathbf{k} \mathbf{E}) + \frac{e \tau_e^2}{m_n^*} \mathbf{k} [\mathbf{B} \mathbf{E}] \right\} \mathbf{k} d\tau_{\mathbf{k}}. \quad (6.93)$$

(6.40), (6.41), (6.42), (6.93)

$$\mathbf{J}_n = \frac{e^{2n}}{m_n^*} \mathbf{E} \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e e^{-\alpha} \alpha^{3/2} d\alpha + \frac{e^3 n}{m_n^{*2}} [\mathbf{B} \mathbf{E}] \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e^2 e^{-\alpha} \alpha^{3/2} d\alpha. \quad (6.94)$$

comes to the view. Average relaxation time for electrons

$$\langle \tau_e \rangle = \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e e^{-\alpha} \alpha^{3/2} d\alpha \quad (6.95)$$

Average of relaxation time squared

$$\langle \tau_e^2 \rangle = \frac{4}{3\sqrt{\pi}} \int_0^\infty \tau_e^2 e^{-\alpha} \alpha^{3/2} d\alpha \quad (6.96)$$

for the circuit gauge of current density, taking into account that $\mu_n = e \langle \tau_e \rangle / m_n^*$

$$\mathbf{J}_n = cn\mu_n \mathbf{E} + \frac{\langle \tau_e^2 \rangle}{|\langle \tau_e \rangle|^2} en\mu_n^2 [\mathbf{B} \mathbf{E}] = en\mu_n \mathbf{E} + r_n en\mu_n^2 [\mathbf{B} \mathbf{E}], \quad (6.97)$$

Get the expression, here

$$\dot{r}_n = \frac{\langle \tau_e^2 \rangle}{|\langle \tau_e \rangle|^2} \quad (6.98)$$

A similar for the hive organizer of vine density

$$\mathbf{J}_p = ep\mu_p \mathbf{E} - r_p ep\mu_p^2 [\mathbf{B} \mathbf{E}], \quad (6.99)$$

We write expression, here

$$r_p = \frac{\langle \tau_h^2 \rangle}{|\langle \tau_h \rangle|^2}. \quad (6.100)$$

Total current density

$$\mathbf{J} = e(n\mu_n + p\mu_p) \mathbf{E} - r_n e \left(\frac{r_p}{r_n} p\mu_p^2 - n\mu_n^2 \right) [\mathbf{B} \mathbf{E}]. \quad (6.101)$$

The coefficients in (6.101) are dependent on the dissolution mechanism of the charge carriers, which determines the time of relaxation. In the case of elastic splashing, and when there is, for any splashing mechanism r_n and r_z $\tau(\mathbf{k})\mathbf{k}' = \mathbf{k}$ and $E' = E$

$$\frac{r_p}{r_n} = \frac{\langle \tau_h^2 \rangle}{|\langle \tau_h \rangle|^2} \frac{|\langle \tau_e \rangle|^2}{\langle \tau_e^2 \rangle} = 1. \quad (6.102)$$

$r_n = r$ and having regard to (6.102):

$$\mathbf{J} = e(n\mu_n + p\mu_p) \mathbf{E} - re(p\mu_p^2 - n\mu_n^2) [\mathbf{B} \mathbf{E}]. \quad (6.103)$$

Let's assume that the current density is directed along the x-axis, i.e., and , the magnetic field is directed along the z-axis, i.e., (Fig. 6.7). In this (6.103) equality is divided into two: $J_x = J$ and $J_y = J_z = 0 = B_z = B$ is $B_y = B_x = 0$

$$rJ_x = e(n\mu_n + p\mu_p)\mathcal{E}_x + re(p\mu_p^2 - n\mu_n^2)B\mathcal{E}_y \quad (6.104)$$

$$J_y = e(n\mu_n + p\mu_p)\mathcal{E}_y - re(p\mu_p^2 - n\mu_n^2)B\mathcal{E}_x \quad (6.105)$$

This solves the system of equations with respect to the Hall field as a weak magnetic field, so that the Hall field is strengthened $\mathcal{E}_y$

$$\mathcal{E}_y = \frac{r}{e} \frac{p\mu_p^2 - n\mu_n^2}{(n\mu_n + p\mu_p)^2} JB = RJB, \quad (6.106)$$

we get the expression where is the Hall constant for a semiconductor with two different types of charge carriers

$$R = \frac{r}{e} \frac{p\mu_p^2 - n\mu_n^2}{(n\mu_n + p\mu_p)^2}. \quad (6.107)$$

For Private Semiconductor $n = p = n_i$,

$$R_i = \frac{r}{en_i} \frac{\mu_p - \mu_n}{\mu_p + \mu_n} = \frac{r}{en_i} \frac{1 - b}{1 + b}. \quad (6.108)$$

Because it is usually $b = \mu_n/\mu_p > 1$ $R_i < 0$.

As it turns out, the Hall constant depends on the scattering mechanism, which is accounted for by magnitude, which is called the Hall factor. For different decomposition mechanisms we delineate. Relaxation time when scattered in acoustic oscillations of the grid taking into account the following:

$$rrE = kT\alpha$$

$$\tau_l = \frac{\tau_n l}{m_n^{*3/2}} T^{-1} E^{-1/2} = \frac{\tau_0 l}{m_n^{*3/2} k^{1/2} T^{3/2}} \alpha^{-1/2} \quad (6.109)$$

Based on (6.95), (6.96) and (6.98) it will be equal to:

$$r = 3\pi/8 \quad (6.110)$$

Based on (5.70) when scattered in input ions

$$r = 315\pi/512 = 1,93. \quad (6.111)$$

Since the scattering of charge carriers depends on temperature, in determining the Hall coefficient at low temperatures, it is necessary to assume $r = 1.93$, when scattering plays a major role in input ions. For the temperatures at which scattering occurs in acoustic phonons, $r = 3/\pi 8$. In the case in which both lattice oscillations and input ions are simultaneously involved in the scattering process, p has a more complex expression. Fig. 6.8 shows that the specific resistance of r to scattering only in input ions of a semiconductor depends on the ratio of r to the ratio of r .

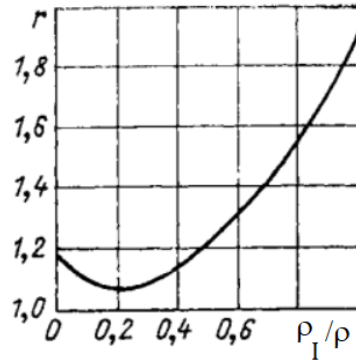


Fig. 6.8. $r = \frac{R}{1/(en)}$ dependence of the coefficient ρ_I/ρ . Donor and Receptor Concentrations, see cm^{-3} ; $1n = 4,3 \cdot 10^{19}$; $2p = 4 \cdot 10^{16}$; $3n = 1,3 \cdot 10^{10}$; $5p = 2,2 \cdot 10^{16}$; $5p = 6 \cdot 10^{18}$; $6n = 1 \cdot 10^{17}$;

For an electron semiconductor with $n \gg p$, based on (6.107), the Hall constant is:

$$R = -r \frac{1}{en} \quad (6.112)$$

At the same time, the specific conductivity of such a semiconductor is determined by the drift mobility and is equal to. So that $\sigma = en\mu_d$

$$\frac{1}{r} |R| \sigma = \mu_d. \quad (6.113)$$

But

$$|R| \sigma = \mu_H \quad (6.114)$$

Hall's mobility can be determined using equality μ_H (6.113) and (6.114)

$$\mu_H = r \mu_d \quad (6.115)$$

that is, Hall's mobility is proportional to the drift mobility.

The relaxation time is in $\tau = \text{the const}$ section, i.e. when it is not dependent on energy

$$r = 1 \quad (6.116)$$

So that

$$\mu_H = \mu_d \quad (6.117)$$

This happens in metals and decomposed semiconductors. Their Hall coefficient R is independent of the scattering mechanism. Since n and p are functions of temperature, R is also dependent on temperature. As an example, Fig. 6.9 shows the temperature-dependent Hall coefficient for n- and p-type indium antimonide samples. As it turns out, in the low-temperature field, the concentration of electrons and spurs remains constant, and this corresponds to the region of decreasing access. In the field of internal electrical conductivity, $\ln|R|$ It is linearly dependent on cross-temperature (curves (1n, 3n and 6n). Because $R > 0$ for the receiving semiconductor,

R decreases as the temperature increases, passes through zero, and becomes negative. The temperature where $R=0$ is called the inversion temperature. From the curves in Figure 6.9 for the $2p$, $4p$, $5p$ and $7p$ samples, it follows that the higher the temperature of transition and the concentration of the internal electrical conductivity, the higher the temperature of inertia.

6.6. Magnitoric Effect

The curvature of the trajectory of the charge carriers in the magnetic field causes a change in the directed current component along the electric field vector, which is equivalent to the change in the resistance of the semiconductor. This phenomenon is called the magnetoresistive effect (Gaussian effect). In the absence of a magnetic field, the current density is determined by the expression of the specific resistance when the direction J coincides with the direction E of the electric field force.

$$\rho = \frac{1}{\sigma} = \frac{\mathcal{E}}{J}. \quad (6.118)$$

In the presence of a magnetic field, when the direction of E does not coincide with the direction of J , its projection is put in its place $\mathcal{E}$

$$\rho_B = \frac{\mathcal{E}_J}{J} = \frac{(\mathcal{E}^J)}{J^2} \quad (6.119)$$

Taking into account equations (6.107), (6.47), and (6.51), it is possible to rewrite the expression for current to a weak magnetic field (6.103)

$$\mathbf{J} = \sigma \mathcal{E} - R \sigma^2 [\mathbf{B} \mathcal{E}]. \quad (6.120)$$

In a linear approximation of the magnetic field $-B^2$ when and taking into account the had $B^2 \mathbf{E}[\mathbf{B} \mathbf{E}] = 0$

$$J^2 = (\sigma \mathcal{E})^2 \quad (6.121)$$

we will have an expression. We put (6.120) and (6.121) in (6.119):

$$\rho_B = \frac{\sigma \mathcal{E}^2}{(\sigma \mathcal{E})^2} = \frac{1}{\sigma} = \rho \quad (6.122)$$

Thus, the semiconductor resistance does not change in a linear approximation in a weak magnetic field.

To determine the specific resistance of a semiconductor with a magnetic field, it is necessary to use quadratic approximation, that is, for a homogeneous semiconductor located in a weak magnetic field and in the absence of a temperature gradient, it is necessary to solve the equation $((\nabla T = 0))$ (6.25): $e \tau_e B / m_n^* \ll 1$ $(\nabla T = 0) \nabla F = 0$

$$\chi_n = -\frac{e\hbar\tau_e}{m_n^*} \left\{ \frac{\mathcal{E} + \frac{e\tau_e}{m_n^*} [\mathbf{B}\mathbf{B}] + \left(\frac{e\tau_e}{m_n^*}\right)^2 (\mathbf{B}\mathbf{G})\mathbf{B}}{1 + \left(\frac{e\tau_e}{m_n^*} B\right)^2} \right\} \quad (6.123)$$

$\frac{e\tau_e}{m_n^*} B \ll 1$ because . As a result $\frac{1}{1 + \left(\frac{e\tau_e}{m_n^*} B\right)^2} \approx 1 - \left(\frac{e\tau_e}{m_n^*} B\right)^2$

$$\chi_n = -\frac{e\hbar\tau_e}{m_n^*} \left[\mathbf{E} + \frac{e\tau_e}{m_n^*} [\mathbf{B}\mathbf{E}] + \left(\frac{e\tau_e}{m_n^*}\right)^2 [\mathbf{B}[\mathbf{B}\mathbf{E}]] \right] \quad (6.124)$$

Current density for a semiconductor with simple-band homogeneous charge carriers

$$\mathbf{J}_n = -\frac{e\hbar^2}{4\pi^3 m_n^{*2}} \int \frac{\partial f_0}{\partial E} \left\{ \tau_e (\mathbf{k}\mathcal{E}) + \left(\frac{e\tau_e}{m_n^*}\right)^2 \mathbf{k}[\mathbf{B}\mathcal{E}] + \frac{e^3 \tau_e^3}{m_n^{*2}} \mathbf{k}[\mathbf{B}[\mathbf{B}\mathbf{E}]] \right\} \mathbf{k} d\tau_{\mathbf{k}}. \quad (5.125)$$

B^2 Integrated in precision up to:

$$\mathbf{J}_n = en\mu_n \left\{ \mathbf{E} + r\mu_n [\mathbf{B}\mathbf{E}] + \frac{3}{2} r\mu_n^2 [\mathbf{B}[\mathbf{B}\mathbf{E}]] \right\} \quad (6.126)$$

We will have it. To solve the resulting equation, $\mathcal{E}$ we spread the vector along the **non-complanar** directions $\mathbf{J}$, $\mathbf{B}$, and $[\mathbf{B}\mathbf{J}]$

$$\dot{\mathcal{E}} = \alpha \mathbf{J} + \beta \mathbf{B} + \gamma [\mathbf{B}\mathbf{J}]. \quad (6.127)$$

We determine the spread coefficients by subtracting (6.126) by (6.127)

$$\mathbf{E} = \frac{1}{en\mu_n} \left\{ \left[1 + \frac{3}{2} r\mu_n^2 \left(1 - \frac{\pi}{4} \right) B^2 \right] \mathbf{J} - \frac{3}{2} r \left(1 - \frac{\pi}{5} \right) \mu\mu_n^2 (\mathbf{J}\mathbf{B})\mathbf{B} - r\mu_n [\mathbf{B}\mathbf{J}] \right\} \quad (6.128)$$

(6,119) by using

$$\frac{\rho_B - \rho}{\rho} = \frac{\Delta\rho}{\rho} = \frac{3}{2} r \left(1 - \frac{\pi}{4} \right) \mu_n^2 \left[B^2 - \frac{(\mathbf{J}\mathbf{B})^2}{j^2} \right] = A\mu_n^2 B^2 \sin^2 \alpha, \quad (6.129)$$

Here the charge depends on the dispersion mechanism of the carriers; - and the angle between $A = \frac{3}{2} r \left(1 - \frac{\pi}{4} \right) \alpha \mathbf{J}\mathbf{B}$:

$$(\mathbf{J}\mathbf{B})^2 = j^2 B^2 \cos^2 \alpha. \quad (6.130)$$

It follows from (6.129) that if $\mathbf{J}$ ($\alpha = 0$) is parallel to $\mathbf{B}$, the resistance of a semiconductor with a spherical symmetry zonal structure does not change, i.e., the longitudinal magnetic resistance is zero. The transverse magnetic resistance ($\mathbf{J} \perp \mathbf{B}$) in weak fields is proportional to the square of magnetic induction and to the square of charge carrier mobility.

To determine the magnetic resistance in a semiconductor with two different types of charge carriers, it is necessary to take a cove organizer of the current density into the hives. (6.126)

$$\mathbf{J}_p = ep\mu_p \left\{ \mathbf{E} - r\mu_p [\mathbf{B}\mathbf{E}] + \frac{3}{2} r\mu_p^2 [\mathbf{B}[\mathbf{B}\mathbf{E}]] \right\}. \quad (6.131)$$

Equation of total current

$$\mathbf{J} = \mathbf{J}_n + \mathbf{J}_p = \sigma \left\{ \mathbf{E} + r \frac{n\mu_n^2 - p\mu_p^2}{n\mu_n + p\mu_p} [\mathbf{B}\mathbf{E}] + \frac{3}{2} r \frac{n\mu_n^3 + p\mu_p^3}{n\mu_n + p\mu_p} [\mathbf{B}[\mathbf{B}\mathbf{E}]] \right\} \quad (6.132)$$

Looks like (6.126) and is solved like (6.126).

The result is for a semiconductor with two different types of conductivity with a simple zone

$$\frac{\Delta\rho}{\rho} = \frac{3}{2} r \frac{e}{\sigma} \left[(n\mu_n^3 + p\mu_p^3) - \frac{\pi (n\mu_n^2 - p\mu_p^2)^2}{4 (n\mu_n + p\mu_p)} \right] B^2 \sin^2 \alpha, \quad (6.133)$$

we will have an expression. It comes to (6.129) when the expression is (or). $n = 0$ or $p = 0$

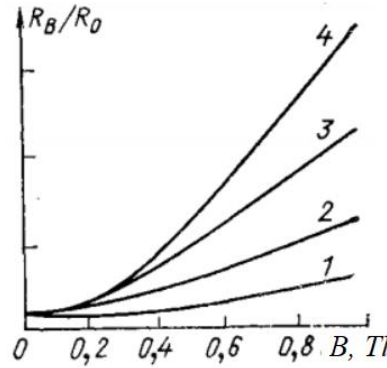


Fig. 6.10. Dependence of the magnetoresistive effect on the magnetic field in indium antimonide

In materials without spherical symmetry of zones, the change in resistance in the magnetic field depends on the direction of the $\mathbf{B}$ vector with respect to the crystallographic axes. In general, the effect of longitudinal magnetic resistance is also observed. Important information about the zonal structure parameters of a given material can be obtained by spinning the specimen on a magnetic field and measuring $\Delta r/r$. The variation in the resistance of a semiconductor wafer in the magnetic field depends significantly on its configuration (the ratio of length to width). The transverse electric field resulting from the Hall effect prevents a change in the trajectory of the charge carriers under the influence of the magnetic field, thereby reducing the magnetic resistance effect. Specifically, it follows from the formula (6.133) that the $\Delta r/r$ ratio is maximum when the second factor is extinguished, which, according to (6.107), occurs when $R=0$, i.e., when there is no Hall field. Given this feature, short, wide specimens with ohmic contacts on their long sides, in order to achieve the greatest change in resistance in the magnetic field, it is necessary to tune the Hall EYuK: The optimal design is a Corbino disk, in which one contact is located in the center of the circular plate and the other is located along

its circumference. In such a sample, the current flows in radial directions, and the Hall EYuK is completely absent. Figure 6.10 shows the change in resistance in the magnetic field for rectangular plates of indium antimony of different configurations, as well as for Corbino disks (R_0 , R_B is the resistance of the sample in the absence and presence of a magnetic field).

Now let us consider the case where there is a strong magnetic field $\frac{e\tau_e}{m_n^*}B \gg$

1. When there is no temperature gradient for a homosexual semiconductor . We extend the vector function by parameter to a line: $(6.25)m_n^*/\ell\tau_e B$

$$\begin{aligned}\chi_n &= -\frac{e\hbar\tau_e}{m_n^*} \left\{ \frac{(\mathbf{B}\mathbf{g})\mathbf{B}}{B^2} + \frac{m_n^*}{e\tau_e B^2} \cdot [B\mathcal{E}] + \left(\mathbf{E} - \frac{(\mathbf{B}\mathbf{g})\mathbf{B}}{B^2} \right) \left(\frac{m_n^*}{e\tau_e B} \right)^2 \right\} = \\ &= -\frac{e\hbar\tau_e}{m_n^*} \left\{ (\mathbf{e}_B \mathcal{E})\mathbf{e}_B + \frac{m_n^*}{e\tau_e B} [\mathbf{e}_B \mathcal{E}] + (\mathcal{E} - (\mathbf{e}_B \mathcal{E})\mathbf{e}_B) \left(\frac{m_n^*}{e\tau_e B} \right)^2 \right\} =\end{aligned}\quad (6.134)$$

where $\mathbf{e}_B$ is the unit vector in the direction of the magnetic field; $-\mathbf{e}_B$ – an organizer of the electric field strength in the direction of the magnetic field; $-\mathbf{e}_B$ – An organizer of the electric field strength in the direction perpendicular to the magnetic field. $\mathbf{e}_B = \mathbf{B}/B$, $\mathbf{E}_B = (\mathbf{e}_B \mathcal{E})\mathbf{e}_B$

Solving the current density equation,

$$\mathbf{J} = en\mu_n \mathbf{E}_B + \frac{en}{B} [\mathbf{e}_B \mathbf{E}] + \frac{4}{3r} \frac{en}{\mu_n B^2} E_\perp. \quad (6.135)$$

We take it. Suppose that the current density is directed along the axis and the magnetic field strength is directed along the x z -axis. Without it

$$\left. \begin{aligned} J_x &= -\frac{en}{B} \mathcal{E}_y + \frac{4}{3r} \frac{en}{\mu_n B^2} \mathcal{E}_x = J \\ J_y &= \frac{en}{B} \mathcal{E}_x + \frac{4}{3r} \frac{en}{\mu_n B^2} \mathcal{E}_y = 0. \end{aligned} \right\} \quad (6.136)$$

This solves the system of equations relative to the $\mathcal{E}_x \mathcal{E}_y$

$$\left. \begin{aligned} \mathcal{E}_x &= \frac{4}{3r} \frac{1}{en\mu_n} J \\ \mathcal{E}_y &= -\frac{1}{en} JB. \end{aligned} \right\} \quad (6.137)$$

We will have it. From here we define the Hall constant and the magnetic resistors

$$R = -\frac{1}{en}; \quad (6.138)$$

$$\rho_B = \frac{(\mathcal{E}\mathbf{J})}{J^2} \frac{4}{3r} \frac{1}{en\mu_n} = \frac{4}{3r} \rho. \quad (6.139)$$

When scattered in the thermal oscillations of the grid . So that $r = 3\pi/8$

$$\Delta\rho/\rho = 4/3r - 1 = 32/9\pi - 1 = 0,13 \quad (6.140)$$

Thus, in a strong magnetic field, the Hall constant does not depend on the scattering mechanism, while the magnetic resistance reaches a critical value. In this case, the magnetic resistance will depend on the scattering mechanism at the saturation point.

Mixed Conductivity Semiconductor Full Current

$$\mathbf{J} = e(n\mu_{it} + p\mu_p)B_B + \frac{e}{B}(n - p)[\mathbf{e}_B \mathbf{E}] + \frac{4}{3r} \frac{e}{B^2} \left(\frac{n}{\mu_n} + \frac{p}{\mu_p} \right). \quad (6.141)$$

This is based on the expression and we define that it is. Then, and by specifying them, $J_x = J$ va $J_y = 0$ $\mathcal{E}_x$ and $\mathcal{E}_y$

$$R = \frac{1}{e(p - n)}; \quad (6.142)$$

$$\rho_B = \frac{4}{3r} \frac{\frac{n}{\mu_n} + \frac{p}{\mu_p}}{e(p - n)}. \quad (6.143)$$

We get expressions.

6.7. Thermoelectric phenomena

Thermoelectric phenomena include Zeebek, Peltje, and Thomson effects.

6.7.1. Zeebek Effective

The Zeebek effect, or thermoelectric effect, is the electric current that occurs in two different semiconductor contacts with different temperatures (Fig. 6.11).

The potential differential that arises depends on the temperature differential and material,

$$\alpha = \frac{d\mathcal{E}}{dT} \quad (6.144)$$

It is characterized by the coefficient, which is called the specific thermo EYuK and is equal in quantity to the EYuK corresponding to the unit temperature. α

If the semiconductor is heated unevenly, the energy and concentration of the charge carriers are large at points where the temperature is high. In a homogeneous semiconductor, respectively, the temperature gradient produces a diffusion flux of charge carriers, producing a gradient of energy and concentration. As a result, an electric current occurs.

The current density at the optional point of the specimen is zero, when the chain is open. This means that the current caused by the temperature gradient is compensated by the current generated in the charge division. Thermo EYuK occurs in the sample.

For a homogeneous unseparated semiconductor with a ferric isoenergetic surface of electrons and hollows, we calculate the thermoenergy using kinetic equations. , where is the width of the forbidden zone. Distribution functions for electrons and hollows for semiconductor insulation and , their concentrations and , because $\left(E = \frac{\hbar k^2}{2m_n^*}\right) \left(E' = -\frac{\hbar k'^2}{2m_p^*}\right)$ uchun sKovaklar uchun $F' = -E_g - FE_g f_0 = e^{-(E-F)/kT} f_{0p} = e^{(E'-F')/kT} = e^{(E'+E_g+F)/kT} n = N_c e^{-(E_g-F)/kT} = N_e e^{F/kT}$ $p = N_v e^{-(F'+E_g)/kT} = N_v e^{-F'/kT}$ bo'ladi $E_g = 0$,

Given that in this case, expressions for and $\mathbf{B} = 0$ $\chi_n(6.25)\chi_p(6.26)$

$$\chi_n = \frac{\hbar \tau_e}{m_n^*} \left[\frac{F - E}{T} \nabla T - \nabla(F - e\varphi) \right]; \quad (6.145)$$

$$\chi_p = -\frac{\hbar \tau_h}{m_p^*} \left[\frac{E' + E_g + F}{T} \nabla T - \nabla(F - e\varphi) \right]. \quad (6.146)$$

will appear to be.

We calculate the current density for the semiconductor, where the charge carriers are scattered in the acoustic oscillations of the grid. In this case, the free-running path would not depend on the energy of the chargers and would be able to express the relaxation time through $l = v\tau l$

$$\tau = \frac{l}{v} = \frac{m_n^* l}{\hbar} k^{-1} \quad (6.147)$$

Considering (6.145), (6.147), and (6.40), the electron organizer of the current (6.32)

$$\mathbf{J}_n = -\frac{c\hbar l}{3\pi^2 m_n^*} \left[\left\{ \frac{F}{T} \nabla T - \nabla(F - e\varphi) \right\} \int_0^\infty \frac{\partial f_0}{\partial E} k^3 dk - \nabla T \int_0^\infty \frac{\partial f_0}{\partial E} E k^3 dk \right] \quad (6.148)$$

can be imagined in appearance. If it is, $E = \alpha k T k^3 dk = \frac{(2m_n^* k T)^2}{2\hbar^4} \alpha d\alpha$

Taking into account (6.6) and (6.145)

$$\mathbf{J}_n = -\frac{4enl}{3(2\pi m_n^* k T)^{1/2}} \left\{ \frac{F}{T} \nabla T - \nabla(F - e\varphi) - 2k \nabla T \right\} \quad (6.149)$$

We take it. Charger carriers using mobility (6.56)

$$\mathbf{J}_n = n\mu_n \left\{ \nabla(F - e\varphi) + \left(2k - \frac{F}{T} \right) \nabla T \right\} \quad (6.150)$$

can be made to look like it. The same is the organizer of the vine for the hives

$$\mathbf{J}_p = p\mu_p \left\{ \nabla(F - e\varphi) - \left(2k - \frac{F + E_g}{T} \right) \nabla T \right\} \quad (6.151)$$

Full Tok

$$\mathbf{J} = n\mu_n \left\{ \nabla(F - e\varphi) + \left(2k - \frac{F}{T} \right) \nabla T \right\} + p\mu_p \left\{ \nabla(F - e\varphi) - \left(2k - \frac{F + E_g}{T} \right) \nabla T \right\} \quad (6.152)$$

It follows from the attitude (6.152) that when a temperature gradient is present in a semiconductor, electromagnetic potential gradient $\nabla(F - e\varphi)$ and temperature gradient ∇T are induced by the action of electromagnetic potential gradient $(F - e\varphi)$ and spitting currents.

To find the thermo EYuK, it is necessary to determine the difference in the potentials for a closed circuit. For an open circuit, if $J=0$ is not, then the voltage of the electric field (6.152) caused by the temperature gradient and called thermoelectric can be found by zeroing the current to zero. However, there are unequal contact potential differences at the limits of semiconductor and measuring metal electrodes due to the present temperature gradient when measuring thermo EYuK (see Item 9.4). The measuring device records the difference in EYuK equal to the thermo EYuK of a semiconductor and the difference in the contact potential of the measuring electrodes. In order to eliminate the last confluent, thermo EYuK must be designated as the electrochemical potential $\nabla\varphi$ gradient. Absolute value $|\nabla F/e|$ Assuming that the metal in the semiconductor and contact will be in thermodynamic equilibrium will be equal to the difference of contact potentials at the boundary points. Therefore, the specific thermo EYuK is determined as follows

$$\alpha = \frac{|\nabla(\varphi - F/e)|}{|\nabla T|}. \quad (6.153)$$

(6.152) by zeroing the current

$$\alpha = -\frac{k}{e} \frac{1}{n\mu_n + p\mu_p} \left\{ n\mu_n \left(2 - \frac{F}{kT} \right) - p\mu_p \left(2 - \frac{F + E_g}{kT} \right) \right\}. \quad (6.154)$$

We take it. If we use the formulas for the concentration of electrons and spikes in an unseparated semiconductor,

$$\alpha = -\frac{k}{e} \left\{ \frac{\left[2 + \ln \frac{N_c}{n} \right] n\mu_n - \left[2 + \ln \frac{N_v}{p} \right] p\mu_p}{n\mu_n + p\mu_p} \right\} \quad (6.155)$$

Here is the density of effective states in the permeability and valent zones. N_c, N_v

Thus, the thermo EYuK of a semiconductor is determined by two conditions, each of which corresponds to the contribution made by electrons and cavities, and these contributions have opposite signs.

From the expression (6.155) it follows that thermo EYuK for an electron semiconductor is negative. In this case, a positive volume charge is generated at the

hot end of the specimen, as the electrons are scattered from the hot end to the cold end. In a hollow semiconductor, the thermo EYuK is positive; In it, on the contrary, a negative volumetric charge appears at the hot end. Thus, if the semiconductor is inletive, then the direction of the internal electric field and the polarity of the thermo EYuK are determined by the gesture of the main charge carriers, and therefore the type of input conductivity of the studied sample can be determined by the thermo EYuK gesture. In the case of mixed conductivity, the thermo EYuK gesture is determined not only by the ratio of concentration of charge carriers, but also by their mobility. In turn, the thermo EYuC of input semiconductors decreases with increasing impurities. When the inlet sample is heated, its thermo EYuK also decreases.

For Private Semiconductor, and (6.154) $\dot{n} = \ddot{p} = n_i, \vec{F} = -\vec{E}_g/2$

$$\alpha = -\frac{k b - 1}{e b + 1} (2 + E_g/2kT) \quad (6.156)$$

takes the view, here $b = \mu_n/\mu_p$.

Thus, the thermo EYuK of a private semiconductor is determined only by the forbidden zone and the ratio of electrons and the mobility of the hives.

The change in the scattering mechanism in a semiconductor, which can occur at different temperatures, affects thermo EYuK through mobility and influences the first wave (within the square bracket) (6.155). For example, its values for acoustic vibrations and input ions in atomic crystals to be scattered by optical vibrations in ion crystals above and below the Debay temperature are 2, respectively; 4; Equal to 3 and 5/2.

Based on the expression (6.155), it is impossible to explain the large increase in thermo EYuK, which is sometimes detected in pure matter at very low temperatures. This phenomenon is defined by the electron-phonon interaction, which is called the phonon takeaway effect.

When a temperature gradient is present in the semiconductor, anisotropy occurs in the scattering of phonons because the concentration gradient of the phonons causes them to move from the hot part of the sample to the cold edge. With each act of electron-phonon scattering that occurs with the absorption of the phonon, the electron acquires the energy and quasi-impulse of the phonon. Since there are more oppositely directed quasi-impulse phonons oriented to the hot-to-cold edge of the sample, the phonon quasi-impulse-absorbing electrons are attracted by the phonon flux. As a result, the electrons begin to slide into the cold part of the sample. The displacement of the electrons continues until an electric field is generated from the stream of the fonon due to their redistribution, which balances the force affecting the electrons. This potential difference represents an additional thermoo EYuK,

which is added to the thermoo EYuK, which is calculated for the absence of electron-phonon interactions. The gravitational effect of the phonon only leads to a significant increase in thermal power at low temperatures, as the phonon distribution gradually returns to equilibrium due to inter-collisions. At high temperatures, the dispersion of phonons on phonons quickly restores the equilibrium distribution of the phonons and the gravitational effect is lost.

The energy relations considered here for the thermo EYuK of semiconductors are in good agreement with the results obtained for germanium at high and low temperatures (Fig. 6.12).

The concentration of electrons in metals does not depend on the haror. Their average energy is:

$$\langle E \rangle = \frac{5}{3} F_0 \left[1 + \frac{5\pi^2}{12} \left(\frac{kT}{F_0} \right)^2 \right] \quad (6.157)$$

is defined by expression, where is the Fermi level at absolute zero. F_0

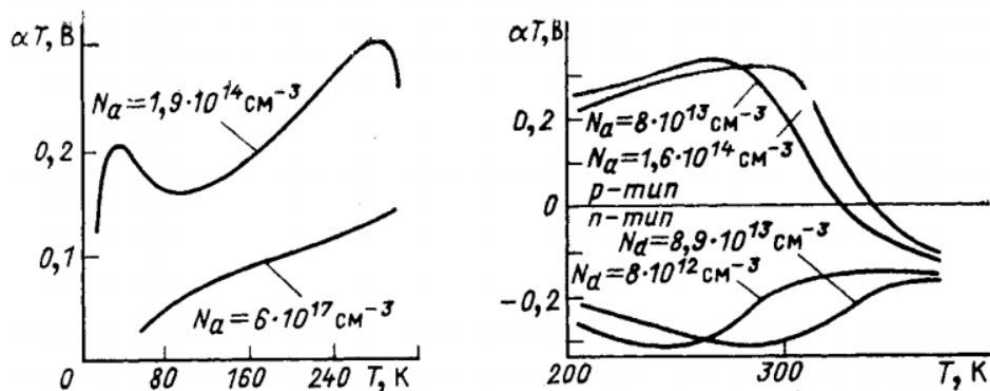


Fig. 6.12. German thermo EYuK at low and high temperatures

Since $F_0 \approx 5 \text{ eV}$ for metals is the second value of (6.157), and the energy of the electrons in the metal is practically independent of temperature. Therefore, the thermo EYuK of metals is smaller than that of semiconductors. For metals, the value reaches $\alpha \approx 1 \div 10 \mu\text{V/K}$ and for semiconductors - $100 \div 1000 \mu\text{V/K}$. Therefore, as a rule, it is not shown to which metal for the thermo EYuK values of semiconductors.

2. Peltie effect

The Peltie effect is the dissipation or absorption of heat transferred by an electric current from two different substance contacts. The discharged heat would be proportional to the current density and time

$$Q_{\Pi} = \Pi j t \quad (6.158)$$

Bu yerda – Peltie koefitsienti. Π

3. Tomson effekti

If there is a temperature gradient along a homogeneous conductor, then heat is emitted (or absorbed) in addition to Joule heat when current passes through the volume of the material, the amount of which is proportional to the amount of electricity flowing and the temperature difference:

$$Q_T = It \int_T^T \tau_T dT \quad (6.159)$$

bu yerda τ_T The coefficient is called the Thomson coefficient.

α , P and the τ_T coefficients are linked to each other through the Thomson relationship. To obtain them, let us consider the thermoelectric effects elementary for the closed circuit of the two conductors, as shown in Fig 6.11. Let the temperature difference between the semiconductor contacts be dT . The presence of a temperature gradient produces a thermoelectric force equal to $dQ = \alpha dT$, which ensures the passage of current dI in this circuit. In this case, the Peltier heat is separated in one of the contacts and absorbed in the other. As a result, Peltier heat is emitted in the chain, and if it differs slightly, then the amount of heat $\Pi' dI - \Pi dI$

$$dQ_{\Pi} = (\Pi' - \Pi)dI = \frac{d\Pi}{dT} dI dT. \quad (6.160)$$

Electric current temperature gradient at 1 s due to the Thomson effect

$$dQ_T = (\tau_{T_2} - \tau_{T_1}) dI dT \quad (6.161)$$

occurs heat dissipation. These thermodynamic processes can be described as reversible since they depend on the current direction and temperature gradient. But in this tight chain there are also irreversible processes – heat conductivity and Joule heat separation. Under stationary conditions, heat conductivity does not change the total amount of heat. Joule heat $(dI)^2$ is proportional to the square of current. Therefore, in the smaller values of dI and dT , they can be ignored. Based on the law of conservation of energy $dQ_{\Pi} + dQ_T = \alpha dI dT$

$$\alpha dI dT = \frac{d\Pi}{dT} dI dT + (\tau_{T_2} - \tau_{T_1}) dI dT \quad (6.162)$$

From here:

$$\alpha = \frac{d\Pi}{dT} + (\tau_{T_2} - \tau_{T_1}) \quad (6.163)$$

According to the second law of thermodynamics if we define it as for the hot end T , for the cold end T' , Π' , τ_{T_2} , Π , τ_{T_1}

$$\frac{\Pi' dI}{T'} - \frac{\Pi dI}{T} + dI \int_T^{T'} \frac{\tau_{T_2} - \tau_{T_1}}{T} dT = 0 \quad (6.164)$$

If we take the value of entropy at its hot end as the numerical head of entropy, we differentiate by (), (6.164) $\Pi' dI/T' =$

$$0T \quad \frac{d}{dT} \left(\frac{\Pi}{T} \right) - \frac{\tau_{T_2} - \tau_{T_1}}{T} = 0 \quad (6.165)$$

we will have an expression. From here

$$\frac{d\Pi}{dT} - (\tau_{T_2} - \tau_{T_1}) = \frac{\Pi}{T}. \quad (6.166)$$

Comparing (6.163) and (6.166)

$$\Pi = \alpha T. \quad (6.167)$$

We clarify that

If we differentiate with respect to (6.167) T

$$\frac{d\alpha}{dT} = \frac{1}{T} \frac{d\Pi}{dT} - \frac{\Pi}{T^2} = \frac{1}{T} \left(\frac{d\Pi}{dT} - \frac{\Pi}{T} \right). \quad (6.168)$$

origin. (6.166)

$$\tau_{T_2} - \tau_{T_1} = \tau_T = T \frac{d\alpha}{dT}. \quad (6.169)$$

Ifodaga ega bo'lamiz.

Thus, the coefficients characterizing the thermoelectric properties of a given substance are linked by the thermodynamic relations between them. Usually measured experimentally, P and P are determined through these relationships. $\alpha, \Pi, \tau_r, \alpha, \tau_T$

6.8. Thermal conductivity of semiconductors

It is known that all bodies conduct heat. This phenomenon is called thermal conduction. If there are no thermal processes leading to the formation of a current in the substance, that is, $J = 0$, then the temperature gradient present in the substance *causes an energy flux W* in the direction opposite to ∇T :

$$\mathbf{W} = -x \nabla T. \quad (6.170)$$

where $x = |W|/|\nabla T|$ Heat is the coefficient of thermal conductivity numerically equal to the amount of energy passed through a single cross-section of the specimen per unit of time, resulting in a temperature difference of one degree at its tips. In a solid, the transfer of thermal energy is carried out by free charge carriers and phonons. The thermal conduction that occurs by the movement of the charge carriers is called electronic (or hives) thermal conduction. It is characterized by thermal conductivity coefficient x_e , which in large part depends on the properties of the semiconductor and temperature. The thermal conduction that occurs by the lattice oscillations is called lattice or phonon thermal conduction. It is denoted by x_L . Thus, the total thermal conductivity of a solid consists of the lattice x_L and the electron (or hollow) x_e thermal conductivity, i.e.

$$x = x'_L + x_e. \quad (6.171)$$

The thermal conductivity of dielectrics is mainly phononic; they have $x_L \gg x_i$. In highly pure metals, the contribution of lattice thermal conductivity is usually small, hence $x_e \gg x_L$. In semiconductors, often lattice thermal conductivity is in the same order as electronic (or hole) thermal conductivity. The lattice thermal conductivity is determined by the flux of phonons passing from the hot end of the sample to the cold end. If we introduce the mean free path l_L for phonons, then the thermal conductivity coefficient of the phonon can be determined on the basis of kinetic theory. In this case, the expression x_L has the following form:

$$x_L = C v_{3I} l_L / 3, \quad (6.172)$$

Here – grill heat sink; – Phonon velocity (sound velocity). $C v_{3B}$

$\mathcal{L}$ In a linear sample in a linear sample, the amount of heat passing through the surface per unit of time when there is a temperature difference between the ends $T_1 - T_2$

$$Q = \frac{x_L}{\mathcal{L}} (T_1 - T_2) \quad (6.173)$$

x_L Using (6.172) to:

$$Q = C (T_1 - T_2) l_L v_{35} / 3 \mathcal{L}. \quad (6.174)$$

We take it. For an unreleased donor semiconductor where charge carriers are scattered in acoustic oscillations, we determine the thermal conductivity caused by the motion of the charge carriers. To do this, we calculate the intensity of the energy flux transported by electrons:

$$\mathbf{W} = \frac{1}{4\pi^3} \int_{(V_E)} f_1^{(n)}(\mathbf{k}) (E - e_\varphi) \mathbf{v}_n d\tau_{\mathbf{k}} =$$

$$\mathbf{J}_n = 0 \quad (6.176)$$

Considering. and we calculate (6.25) given (6.175) for the case $\mathbf{B} = 0$ $\mathbf{J}_n = 0$

$$\mathbf{W} = -n\mu_n \frac{2k^2 T}{e} \nabla T. \quad (6.177)$$

The multiplication in front of $\mathbf{W}$ is the coefficient of thermal conductivity: ∇T

$$x_e = 2k^2 T n \mu_n / e = 2(k/e)^2 \sigma T. \quad (6.178)$$

Consequently, the thermal conductivity of a substance is determined by its temperature and specific conductivity.

6.178 dan

$$x_t / \sigma = 2(k/e)^2 T, \quad (6.179)$$

We get the attitude. This is Wideman-Franz law.

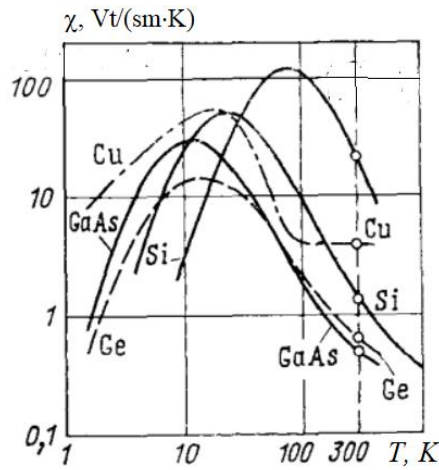


Fig. 6.13. For germanium, silicon, gallium arsenide, copper and diamond, the temperature dependence of conductivity

For thermal conductivity, performing similar calculations for an unsplit semiconductor with two different types of conductivity

$$\chi_e = \frac{2k^2T}{e} n\mu_n + \frac{2k^2T}{e} p\mu_p + \frac{n\mu_n p\mu_p (E_g + 4kT)^2}{eT(n\mu_n + p\mu_p)} \quad (6.180)$$

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As can be seen from Fig. 6.13, the thermal conductivity of solids first increases with increasing temperature, and then decreases at higher temperatures.

6.9. Electrical conductivity of semiconductors in a strong electric field

It has been assumed that the relaxation time in deriving the kinetic equation does not depend on the electric field. Then, using a kinetic equation for an unseparated semiconductor with a spherical isoenergetic surface, it was shown that the current density is subject to Ohm's law, that is, it is proportional to the electric field strength

$$\mathbf{J} = en\mu_0 \mathbf{E} = \sigma_0 \mathbf{E} \quad (6.181)$$

Where μ_0, σ_0 - mobility and specific conductivity in a weak electric field.

When an electron collides with an acoustic phonon, it can transfer energy into the phonon or receive energy from the phonon. Let the electron energy change ΔE . If we calculate the number of collisions associated with the transfer of energy to the background from the total number of collisions, then according to (5.99), for one of these collisions there are $2N_q + 1$ collisions. Hence the relative energy loss of the circuit during the collision

$$\delta = \frac{\Delta E}{kT} = \frac{\hbar\omega_q}{kT} \frac{1}{2N_q + 1} = \frac{\hbar\omega_q}{kT} \text{th} \frac{\hbar\omega_q}{2kT} \quad (6.182)$$

Respectively, the change in the energy of an electron in collision with a phonon. If the energy gained by an electron in an electric field on a free-treaded

path is much smaller than the energy it loses in a single collision with a grid $\Delta E = kT\delta eFl$

$$\frac{e\mathcal{E}l}{\Delta E} \ll 1$$

Collision is called enormous. That's Ohm's Law

$$e\mathcal{E}l \ll kT\delta, \quad (6.183)$$

or

$$\mathcal{E} \ll \delta \frac{kT}{el} \quad (6.184)$$

when it happens.

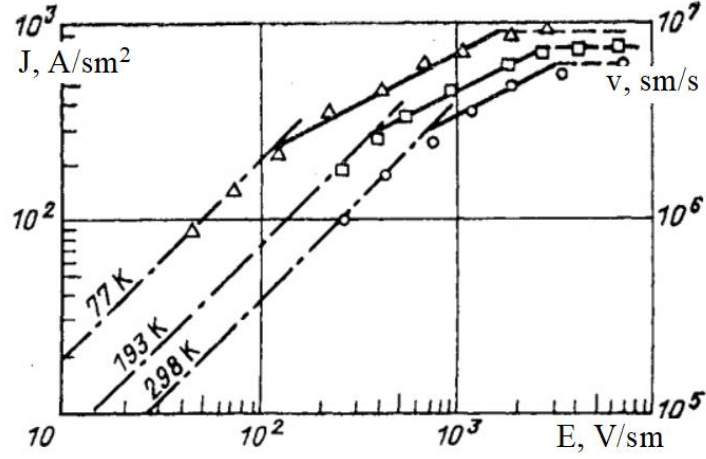
In this case, the distribution of the velocity of the charge carriers becomes an almost Maxwell distribution with a temperature equal to the lattice temperature. In this case, the chaotic motion of the charge carriers is added to their slow displacement caused by the action of their weak external electric field. In this case, the displacement velocity of the charge carriers v will be significantly less than *their heat velocity* v_0 , i.e.

$$v \ll v_0 \quad (6.185)$$

At strong electric fields, Ohm's law is broken. This can happen for two reasons. First, in strong electric fields, the displacement velocity of the charge carriers is proportional to their thermal velocity. It affects the process of digestion, which is manifested in a change in the timing of relaxation and, consequently, motility. This includes the "heating" phenomenon of electron-comb gas and the Gann effect. Second, strong electric fields can cause a variation in the concentration of the charge carriers. In this case, an increase in the concentration of charge carriers may occur for a number of reasons. Such effects are the shock ionization, the Zener effect, and electrostatic ionization.

Consider the electrical conductivity of a semiconductor in a strong electric field. (6.184) The condition is that in strong electric fields where the charge carriers collide with acoustic phonons, they do not manage to give the crystal lattice the energy it receives in the field. As a result, the average energy of the charge carriers increases, and the electron pit gas "heats up." But as the energy of the charge carriers increases, their velocity increases, which reduces the time of free running and, accordingly, increases the likelihood that the charge carriers will scatter into the thermal vibrations of the lattice. Therefore, the increase in charge carrier energy as the field strength increases will continue until the electron-phonon interaction is so strengthened that a steady state is established in the electric field in which the energy received by the charge carrier in the unit of time eCv is equal to the energy lost by Rice

$$e\mathcal{O}v = \delta kT/\tau. \quad (6.186)$$



6.14. Dependence of current density and drift velocity in electron germanium on external electric field strength

Based on the lattice's thermal oscillations (5.93), the free-running track is not dependent on energy, hence the time of the free run. But based on the speed of chargers (6.186): $\tau = l/v_0$

$$v = \delta kT/e\mathcal{E}\tau = \delta kTv_0/e\mathcal{E}l. \quad (6.187)$$

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$$v = \frac{e\tau}{m^*} \mathcal{E} = \frac{el}{m^*v_0} \mathcal{E} \quad (6.188)$$

Comparing (6.187) and (6.188)

$$(e\mathcal{E}l)^2 = 28kTm^*v_0^2/2 \quad (6.189)$$

origin. But since it is, we determine the critical strength of the field from (6.189) $m^* v_0^2 / 2 = kT$

$$\mathcal{X}_{up} = \frac{kT}{ll} (2\delta)^{1/2}, \quad (6.190)$$

At the same time, heating of electron-pit gas begins. It follows from (6.190) that the critical field strength decreases as the temperature decreases. This is consistent with the results of the experiment shown in Fig. 6.14.

Using (6.187) and (6.190),

$$v = (\delta/2)^{1/2} v_0. \quad (6.191)$$

We take it. It follows that the heating of the electron-furnace gas begins when the velocity of the charge carriers in the electric field approaches the heat velocity.

From (6.189), the energy of the charge carriers can be expressed in terms of the heat velocity

$$\left(\frac{m^*v_0^2}{2} \right)^2 = \frac{(e\mathcal{F}l)^2}{2\delta} \quad (6.192)$$

From here

$$v_0 = \left[\frac{2e\mathcal{E}l}{m^*(2\delta)^{1/2}} \right]^{1/2}. \quad (6.193)$$

(6.193) and (6.191) Qoyamiz

$$y = \left(\frac{\delta}{2} \right)^{1/2} \left[\frac{2e\mathcal{E}l}{m^*(2\delta)^{1/2}} \right]^{1/2} = \left[\frac{el}{m^*} \left(\frac{\delta}{2} \right)^{1/2} \right]^{1/2} \mathcal{E} \mathcal{E}^{1/2}. \quad (6.194)$$

Thus, the electric field amplification occurs in strong fields where the critical value is greater than the critical value, i.e.

$$\mathcal{E} > \mathcal{E}_{\text{HP}} \text{ или } \mathcal{E} > \frac{kT}{el} (2\delta)^{1/2} \quad (6.195)$$

the drift velocity of the charge carriers when

$$v \sim \mathcal{E}^{1/2} \quad (6.196)$$

while the relaxation time decreases accordingly

$$\tau = \frac{l}{v_0 + v} \sim \mathcal{E}^{-1/2} \quad (6.197)$$

From here

$$\mu_n = \frac{E\tau}{m^*} = \mu_0 \mathcal{E} - 1/\ddot{z}, \quad (6.198)$$

origins, as well as

$$\sigma = en\mu_n = \sigma_0 \mathcal{E} - 1/2 \text{ и } J = \sigma \mathcal{E} = \sigma_0 \mathcal{E}^{1/2} \quad (6.199)$$

Consequently, if there is an external electric field force, Ohm's law is not fulfilled. For example, in Germany, deviations from ohmic conductivity at room temperature occur at a field strength of about 700 V/cm (Fig. 6.14). In the field of even more powerful electric fields, the energy of the charge carriers increases so much that they can be scattered by optical vibrations. In this case, when an optical phonon collides, the energy of the carrier changes to a value equal to it, and the energy loss rate increases significantly. As a result, a new stationary state is established, wherein the energy received by the charge carrier in a unit of time in the field is equal to the energy lost by that charge carrier at the same time, i.e. $\mathcal{E} > \mathcal{E}_{\text{HP}}$

$$e\mathcal{E}v = \frac{\Delta E}{\tau}. \quad (6.200)$$

But according to (6.182), the energy of the charge carrier lost in a collision

$$\Delta E = \hbar\omega_0 \text{ th } \frac{\hbar\omega_0}{2kT}, \quad (6.201)$$

$$e\mathcal{E}v = \frac{v_0}{l} \hbar\omega_0 \text{ th } \frac{\hbar\omega_0}{2kT} \quad (6.202)$$

If we find the electric field strength from (6.188) and put it in (6.202), the drift velocity of the charge carriers follows

$$v = \left(\frac{\hbar\omega_0}{m^*} \text{th} \frac{\hbar\omega_0}{2kT} \right)^{1/2} \quad (6.203)$$

Its marginal value

$$v_s = (\hbar\omega_0/m^*)^{1/2} \quad (6.204)$$

Thus, the density of the vine passing through the sample tends to slide and

$$J_s = env_s = en(\hbar\omega_s/m^*)^{1/2} \quad (6.205)$$

Not at all. Mum

$$\sigma = \frac{J_c}{\mathcal{E}} = \sigma_0 \mathcal{C}^{-1}. \quad (6.206)$$

signifies that there is a lot of work to be done.

Electric field strength from (6.202) and (6.203)

$$\mathcal{E}'_{kp} = \frac{v_0}{el} \left(m^* \hbar \left(\omega_0 \text{th} \frac{\hbar\omega_0}{2kT} \right)^{1/2} \right) \quad (6.207)$$

In covalent crystals, saturation of current and drift velocity is observed.

The current and drift velocities will not depend on the external electric field. Estimates indicate that saturation of current and dreuff velocity $\mathcal{E} > \mathcal{E}'_{kp} \approx 10^4$ V/sm occurs when there is field strength , which reaches drift velocity. $\approx 10^7$ sm/s

This is consistent with the experimental data in Fig. 6.15, which shows that electrons and velocities for pure germanium and silicon depend on the electric field. From the information in this figure, it follows that in weak fields, Ohm's law is satisfied. Starting from the critical value of field strength that depends on temperature, it is proportional to the current. In very strong fields, the current will not depend on the field. $\mathcal{E}^{1/2}$

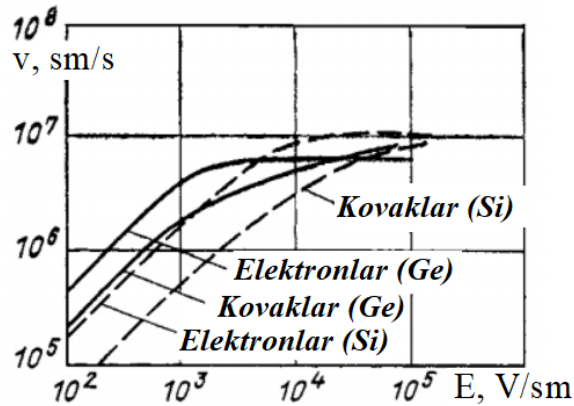


Fig. 6.15. The dependence of the velocity of the charge carriers on the electric field strength for pure germanium and silicon

In polar crystals, scattering in optical oscillations is also prevailing at low values of the energy of the charge carrier. In external electric fields with electron

energy, mobility decreases with field amplification, as in covalent semiconductors. When very strong fields have electron energy, instead of decreasing the mobility and saturation of the shear velocity, their rapid growth occurs and the energy of the charge carriers increases. This phenomenon is called the "escape" of chargers. $E \ll \hbar\omega_0$ $E \gg \hbar\omega_0$

6.10. Gann effekti

In semiconductors where the conductivity band has more than one energy minimum, when an electron is scattered in which the wave vector k corresponds to one of the miniatures, the wave vector belonging to the other minimum can fall into the state k' . As a result of such scattering, the transition of electrons from one minimum of the conduction zone to another occurs. This type of sweeping is called inter-valley.

Under certain conditions, inter-valley scattering of charge carriers can cause the occurrence of frequency current oscillations of the order of 10^{10} Hz when a strongly immutable electric field is applied to a single-species semiconductor. This phenomenon, called the Gann effect, was first observed in the Gaas.

Figure 6.16 shows the energy structure of gallium arsenide in the direction of the [100] axis. What is important here is the presence of minimas A and B, separated by a gap of $\Delta E = 0.36$ eV, where the effective masses of the electrons are different. In the Lower A valley region, electrons have a light, effective mass $m_1^* = 0,068m_0$ and high mobility $[\mu_1 \approx 4000 + 8000 \text{ sm}^2/(V \times s)]$. In the region of the upper B valley, electrons are heavy and have low mobility with $m_2^* = 1,2m_0$ $[\mu_n \approx 100 \text{ sm}^2/(V \cdot s)]$. The density of cases is 70 times higher in the upper part than in the bottom.

At sufficiently low temperatures and with a weak external electric field, the electron energy E will be significantly less than the energy gap $kT \ll \Delta E$, i.e., $E \ll \Delta E$, so the electrons occupy energy levels mainly in the lower valley of the conduction zone (Figure 6.17, a), where they have a high mobility and low mass state. In this case, the current density is determined by the concentration of electrons $n_1 = n_0$ in the lower valley ($n_2 = 0$ in the upper valley), their mobility is determined by μ_1 , and increases linearly with an increase in field strength up to a certain critical value Ea (Fig. 6.18):

$$J = en_0\mu_1\mathcal{E}. \quad (6.208)$$

As the electric field strength increases, the mean energy of the electrons increases, and it becomes possible for electrons to pass into the upper valley (figure 6.17, Fig. b), where they will have small mobility, greater effective mass, and a higher state density, the resulting concentration of electrons. $E > \Delta E$ $n_0 = n_1 + n_2$

Thus, as the field strength increases, the mobility of the electrons and the current density decrease, a cumulative field in the volt-ampere characteristic occurs.

$\mathcal{E} > \mathcal{E}_\zeta$ when all electrons are in the upper valley and the current density is

$$J = en_0\mu_2\mathcal{E} \quad (6.209)$$

will be equal to .

We will consider such a mechanism of frustration. Suppose an external voltage is placed on a specimen of length. In a homogeneous semiconductor sample, the electric field over its entire length is the same. But if the sample has a dissonance with a locally high resistance, the electric field strength here is much higher. Accordingly, the field strength at this point reaches a critical value. $\mathcal{L}\mathcal{E}_{kp} = \mathcal{E}_a$

As soon as the field strength in the local irregularity region $\mathcal{E}$ reaches critical a the electrons pass into the upper B valley, and heavy electrons appear in this region of the sample where the field is strengthened. The mobility of the electrons in this part of the sample decreases and its resistance increases. This leads to an increase in field strength at this point in the sample, which causes a more intense transition of electrons into the upper valley. However, since the voltage applied to the sample does not change, the area tension to the right and left of this region of the sample decreases. As a result, the distribution of the electric field is not dramatically uniform, and a region of a strong electric field is formed, which is called the electric domain (Fig. 6.19, a).

Under the influence of the electric field of the domain, the region of heavy electrons moves across the sample at a relatively low speed due to the low mobility of the heavy electrons. On the right and left sides of the heavy circuit zone, light electrons move at high speeds. On the left side, they reach this zone, and as a result an area of increase in electron concentration is formed - a negative spatial charge area. On the right side of this zone, the electrons of light move forward, so a region is formed in which the electron is reduced - the positive spatial charge region. Consequently, in the strong electric field region of the domain, on the circuit concentration distribution curve there is a diminished region of $n < n_0$ corresponding to the domain façade and a region enriched with $n > n_0$ corresponding to domain decline (Fig. 6.19, Fig. b).

$\mathcal{E}_d$ The velocity of electron movement $v_{1\mu}$ also increases as field strengths are increased within the domain . The field strength outside the domain decreases dramatically, so the speed of movement of electrons decreases. The increase in field strength in the domain $\mathcal{E}_v$ stops when the velocity of the domain v_d is equal to the displacement velocity of the electrons outside the domain v_n , i.e.

$$v_d = v_n \quad (6.210)$$

or

$$\mu_2 \mathcal{E}_\mu = \mu_1 \mathcal{E}_B \quad (6.211)$$

In this case, the fixed velocity of the vd domain will be less than the maximum velocity that the electrons have. Therefore, when the voltage is connected to the sample $\mathcal{E}_a$ the $t=t_1$ current will have the maximum I_{max} value determined by v_{max} . The domain formation process starts immediately and since this process is short-lived, since the time constant associated with the intermittent passage of the electron is equal to s , the current very 10^{-12} quickly falls to the I_{min} value determined by the speed of movement of the domain v_2 at steady state:

$$I_{min} = se n_0 v_\mu, \quad (6.212)$$

Here is the cross-sectional surface of the sample. s

The minimum current value is retained for the entire time of domain movement over the sample. As marked, this time is determined by the length of the sample and the speed at which the domain travels:

$$T = \frac{\mathcal{L}}{v_\mu}. \quad (6.213)$$

The domain goes from the cathode to the anode. When it reaches the anode, the strong field region leaves the specimen, and the current in the sample begins to intensify. When the current in the sample reaches the I_{max} value, a new domain is formed and the thread falls to *the* I_{max} . As a result of the movement of the domain around the crystal, a current impulse occurs in the outer chain. The work cycle of the current impulses is determined by the time T passes through the domain. For a $50 \mu m$ long sample, the current oscillation frequency should be about 2 GGs. An approximate view of these oscillations is shown in Fig. 6.20.

Since the domain formation effect is a volumetric instability, it does not depend on the external elements of the scheme. This means that the rate of the domain does not depend on the magnitude of the outside voltage and the domain moves around at a drift speed of about 10^7 cm/s. While there may be a few inconsequentialities from which domains can appear in a crystal, there is only one domain in a crystal. In other words, the appearance of a domain only happens in one of the homosexuality. Once the domain is lost, a new domain with a different gender identity can appear. The different arrangement of homogeneities in the crystal determines the different transition times of the domains, i.e., different oscillation periods. Therefore, pure and very homogeneous samples are required to observe the Gann effect. Due to the inhomogeneity of the degree of alloying that occurs during the production of electrodes in plate samples, domains are formed in the region of the magnified area near the cathode. It is also important that the distance between the minimum A and B of the conduction band is not very large, since the electron's

transition into the second valley requires a high intensity area, in which an increase in electron concentration is possible due to shock ionization or a tunneling effect.

6.11. Zarbli ionizatsiya

In a strong electric field, if it increases to about 10 V/cm, the electrons (or holes) receive enough energy to ionize the atoms of the substance. As a result, electron-hole pairs emerge (Fig. 6.21), which in turn are accelerated by the field and form additional charge carriers. The process of creating charge carriers is called shock ionization. However, this phenomenon can only be observed in $p-n$ transitions, since it is very difficult to create V/cm fields at the optimal thickness of the other in a homogeneous semiconductor. Except for the shock ionization of the input atoms, which requires small areas and is observed at low temperatures when the input atoms are not ionized. The fifth and third group elements - germanium, shock ionization of donors and receptors occur in areas equal to approximately 5-10 V/cm. For example, from Figure 6.22, which shows the dependence of the current density on the electric field in the electron germanium at $4.2 \cdot 10^5 - 10^6$ K, it follows that there is a sharp jump-like increase in current density caused by the ionization of the input atoms at $E \approx 5$ V/cm. Rate of formation of electron-hole pairs due to shock ionization

$$G = \alpha_n n \mu_n + \alpha_p p \mu_p \quad (6.214)$$

given in the form, in which α_n and α_p are the shock ionization coefficients of electrons and holes, respectively.

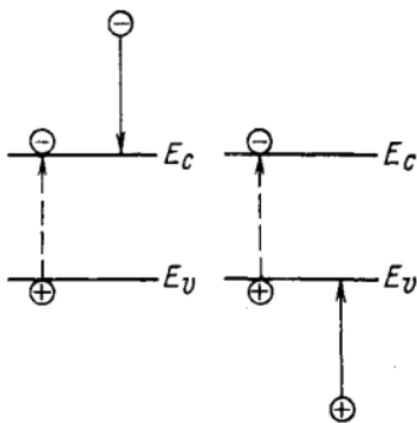


Fig. 6.21. Two Different Variants of Shock Ionization

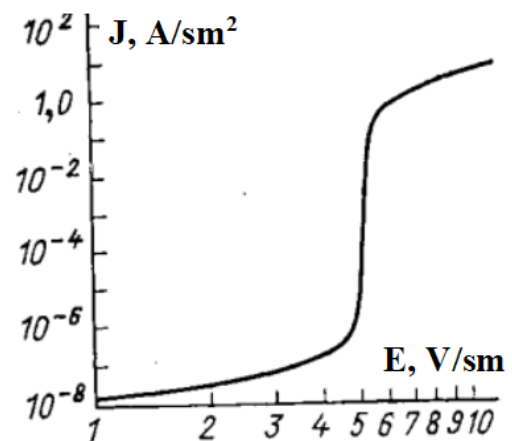


Fig. 6.22. Shock ionization in n-germanium at $T = 4.2$ K

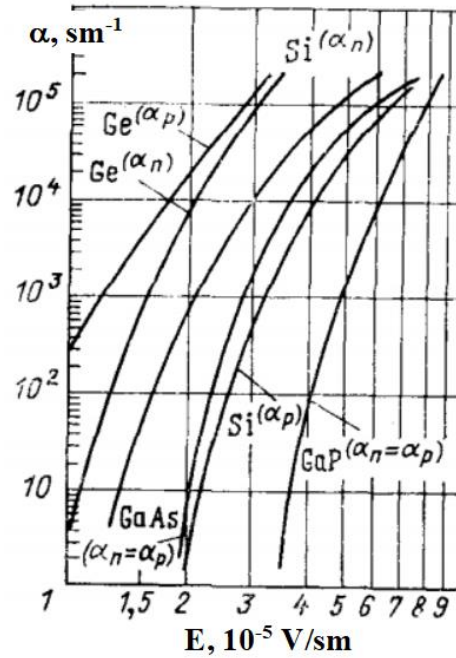


Fig. 6.23. Dependence of the charge ionization coefficient for germanium, silicon, gallium arsenide, and gallium phosphide on the electric field

It is defined as the number of electron-spur pairs created by an electron (spit) per unit escape distance. The shock ionization coefficient of both electrons and their spurs is strongly dependent on the value of the electric field strength. Fig. 6.23 shows the experimentally measured values of shock ionization coefficients for germanium, silicon, gallium arsenide, and gallium phosphide. $\alpha_n \alpha_p$

Shock ionization coefficient $\alpha(E)$

$$\alpha(E) = \frac{1}{v} \int_0^\infty W_n(E) f(E) N(E) dE \quad (6.215)$$

is determined by the relation where is the probability of ionization, calculated by quantum mechanics; – a relative distribution function; – the density of quantum states; – drift speed. $W_n(E) f(E) N(E) v$

Ionization probability (Ge and Si) for semiconductors with high dielectric absorption :

$$W_u(E) \approx A(E - E_u)^2, \quad (6.216)$$

where – the coefficient of proportionality, is the ionization energy. AE_u

Charge ionization involves three charge carriers that distribute the energy of the initial charge carrier approximately equally; boundary ionization energy, where $EE_u = \frac{3}{2} E_{gg}$ is the forbidden zone width. The α field dependence of the coefficient of exposure (E) is largely determined by the type of the distribution function $f(E)$. A detailed analysis using the kinetic equation shows that it is not possible to obtain a

full analytic expression for the distribution function in the $0 < E < \infty$ energy domain; $f(E)$ can be obtained explicitly for two spheres: $E < E_i$ and $E > E_i$ spheres.

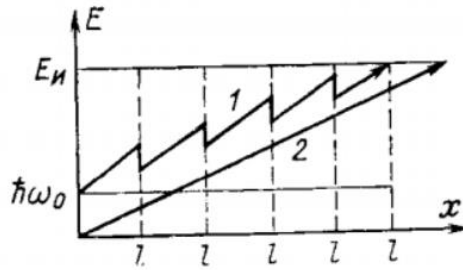


Fig. 6.24. Shock ionization conditional diagram

The acceleration of E_i to the ionization limit of electrons depends on the ratio of two factors - acceleration in the outer electric field and the dissipation of energy during collisions with phonons. The greatest contribution to the energy release is made by optical phonons. The energy lost by an electron in a collision with an optical "phonon" is equal to $\hbar\omega_0 = \text{const}$ and the mean free path is not dependent on energy, i.e., $l = \text{const}$. $E_p \hbar\omega_0$ The acceleration to the ionization boundary can be shown by the following conditional diagram, where the electron shown in Fig. 6.24 is drawn vertically and the average distance of the electron between the two ionizing collisions is drawn horizontally. s with phonons, since the relative energy in each such collision is small. The loss is small. In this case, the theory shows that,

$$\alpha(E) \sim e^{-E_H/\mathcal{E}^2}. \quad (6.217)$$

A similarity with the diffusion action of a particle can be applied here, so the α approximation denoted for $\alpha(E)$ is called diffusion.

In the second case, when the fields are relatively weak ($eEl \ll \hbar\omega_0$), the charge carriers receive energy at a distance equal to many average free paths (Row 2). In such a situation, there is only one chance of receiving E_r energy when the charge carriers do not collide with phonons before the ionization boundary. Since the path to the ionization boundary and the mean free path between collisions with phonons is $\mathcal{L}_n = E_n/e\mathcal{E}l$, there is a probability of going to the ionization boundary without collisions. In this connection, the distribution function and therefore the ionization coefficient of the effect $e^{-\mathcal{L}_n/l} \approx e^{-E_n/\rho_n Z t}$

$$\alpha(E) \sim e^{-E_u/e\mathcal{E}l}. \quad (6.218)$$

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The distribution function is strongly extended along the direction of the electric field, that is, in the direction of the displacement of the carrier, so this approximation is called drift.

As experiments have shown, in Ge and Si (E) it is largely determined by the law of (6.218). However, in narrow p-n transitions (this is especially evident in

Si) (6.217), an apparent correlation is observed. This occurs when $E \geq 5 \cdot 10^5$ V/cm field is amplified.

6.12. Tunneling effect and electrostatic ionization

In very strong electric fields, another mechanism of formation of free charge carriers will be possible: the tunneling effect or Ziner effect. It is known that a semiconductor located in an external electric field has the bending of energy zones (see Section 2.10). The greater the electric field strength, the greater is the slope of the zones. Fig. 6.25 shows the zonal structure of a private semiconductor with a strong electric field. In this case, the transition of the electron from the stress zone to the permeability zone is possible due to the tunneling effect. The likelihood of going into a tunnel depends on the height and width of the potential barrier.

In this case, the height of the potential barrier AB is the width of the prohibited E_g zone. The effective barrier width $AB = \Delta x$ can be determined from the difference in the potential energy of the electron in the conduction zone at point B and the potential energy of the electron in the tensile zone at point A. Because of the potential energy up to the extra constant.

$$U(B) - U(A) = -e\mathcal{E}\Delta x = -E_g \quad (6.219)$$

Effective width of barrier

$$\Delta x = \frac{E_g}{e\mathcal{E}}. \quad (6.220)$$

So the potential barrier width depends on the strength of the electric field.

The transition of an electron from point A to point V is due to the passage of ABV through a triangular barrier. But, as quantum mechanics knows, the probability of passing a triangular barrier

$$D = D_0 e^{-\frac{8\pi\sqrt{2m_n^*}E_g^{3/2}}{3eh\mathcal{E}}}. \quad (6.221)$$

looks.

It follows from this expression that the probability of a tunnel going from one zone to another increases exponentially with the increase in the electric field E . In addition, the probability of tunneling is higher for semiconductors with small values and effective mass m^* of the E_g banned band (Fig. 6.26).

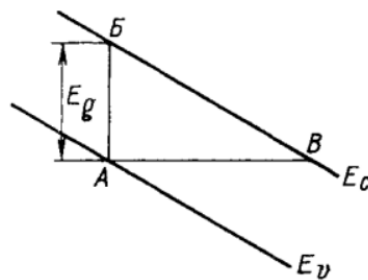


Fig. 6.25. Zonal structure of a private semiconductor under a strong electric field

For transitions from the tensile zone to the permeability zone, as well as from the permeability zone to the stress zone, the probability of tunneling will be the same. But since there are more electrons in the stress band than in the conduction zone, a preferential transition of electrons from the stress band to the conduction zone occurs. Therefore, the tunneling effect can lead to a significant increase in the concentration of free charge carriers. It is observed in areas of 10^6 V/cm.

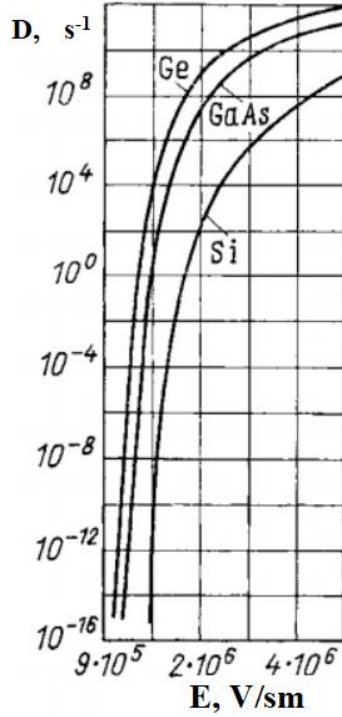


Fig. 6.26. Dependence of the probability of tunneling an electron in a p-n-transition on the electric field

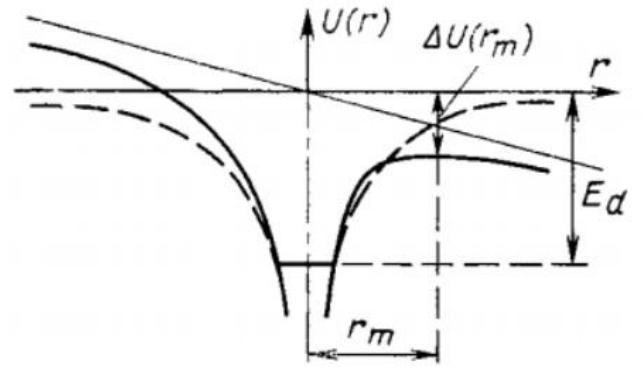


Fig. 6.27. Drop of potential barrier in external electric field

The increase in the concentration of free charge carriers may also be associated with a decrease in the ionization energy of the donor input atoms in a strong electric field. Fig. 6.27 shows the potential energy of the electron in the donor input atom with no field (linear curve) and external electric field (continuous curve). In the absence of an electric field, the electron is localized in the potential well generated by the positively charged donor, and the energy required for its ionization is equal to E_d .

The external electric field lowers the potential gradient in the direction opposite to the direction of the electric field. Taking into account the effect of the external electric field directed along the vector $\mathbf{p}$, the potential energy of an electron can be written as:

$$U(\mathbf{r}) = -e\mathcal{E}r - \frac{e^2}{4\pi\epsilon_r\epsilon_0 r} \quad (6.222)$$

The first state in this equation refers to the change of potential energy in the external electric field, and the second in the donor's coulomb field. It follows from (6.222) that at some point one reaches the extreme, the potential obstacle at this point $|U(\mathbf{r})|$ is r_m

$$\Delta U(\mathbf{r}_m) = -e\sqrt{e\mathcal{E}/\pi\epsilon_r\epsilon_0}, \quad (6.223)$$

decreases. Consequently, the ionization energy of the donor is also reduced by this amount.

The decrease in the ionization energy of the donor is based on Boltzman's statistics

$$e^{-\Delta U/kT} = e^{\frac{(e^3/\pi\epsilon_r\epsilon_0)^{1/2}}{kT}\sqrt{\delta}} \quad (6.224)$$

ga oshiradi.

The electric field strength, where a significant increase in electron concentration occurs, is typically $10^4 - 10^5$ V/cm. This mechanism of occurrence of free charge carriers manifests itself only when the concentration of non-ionized donor input is significantly exceeded by the concentration of free electrons.

Control Questions

1. What is a distribution function?
2. How is the specific electrical conductivity of semiconductors determined?
3. What is the mobility of charge carriers?
4. What is the Hall effect?
5. How does the Hall effect occur in semiconductors with two types of charge carriers?
6. What is the magnetoresistive effect?
7. Explain thermoelectric phenomena?
8. How is the thermal conductivity of semiconductors determined?
9. What is the Gann effect?
10. Ezplane Tunnel effekt.

Chapter VII.

GENERATION AND RECOMBINATION OF ELECTRONS AND HIVES

7.1. Balanced and non-equilibrium charge carriers

Consider a donor semiconductor that is in a state of thermodynamic equilibrium. As a result of thermal generation, electrons from the donor mixture pass into the conduction zone. At sufficiently high temperatures, the transition of electrons directly from the stress zone to the conduction zone becomes noticeable. Fig 7.1 shows the passage of electrons through the axes as a result of thermal excitation. Free charge carriers formed as a result of thermal generation and being in thermodynamic equilibrium with a crystal lattice are called equilibrium charge carriers.

Fig 7.1 shows the density of the quantum states in the conductivity and valent band $N(E)$, the Fermi-Dirac distribution function $f_0(E)$, and the concentrations of the electrons n_0 and p_0 of the swells in equilibrium that occupy states near the edges of the respective zones (barred areas).

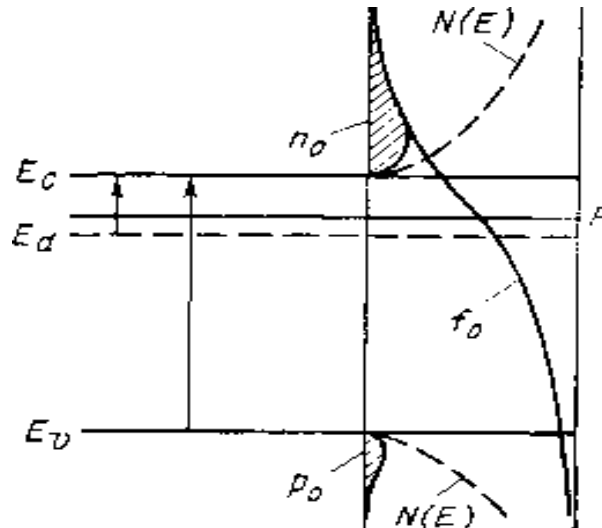


Fig. 7.1. Thermal generation of free charge carriers and their concentration at the thermodynamic equilibrium state

With the generation of free charge carriers, a recombination process occurs simultaneously: electrons return to the free states in the stress zone, as a result of which the free electron and the free spark are lost. Under the conditions of thermodynamic equilibrium, the processes of generation and recombination are mutually equilibrium. We define the number of generated electronic hives pairs per

crystal volume unit by G_0 and the number of recombinable electronic pairs by R_0 . The probability of recombination is proportional to the multiplication of the concentration of free charge carriers, so

$$R_0 = \gamma_r n_0 p_0 \quad (7.1)$$

where γ_r is the coefficient of proportionality, which is called the recombination coefficient.

For the equilibrium state of a semiconductor, parity applies

$$G_0 = R_0 \quad (7.2)$$

This is an expression of the principle of elaborate balance.

Besides thermal generation, there are other mechanisms which lead to the emergence of free charge carriers. For example, free charge carriers can also be generated by semiconductor illumination, tensile bonds breaking at strong electric fields, or by injection using p-n transition. In all of these cases, known concentrations of electrons n and p of electrons occur, which differ from those in thermodynamic equilibrium of free charge carriers.

Moving charge carriers, which are not thermodynamically equilibrium both in concentration and in energy distribution, are unequilibrium charge carriers, and their concentration n, p is called unequilibrium. The difference between the non-equilibrium concentration of charge carriers n, p and the concentration of n_0, p_0 in equilibrium $\Delta n, \Delta p$ is commonly referred to as the excess concentration of charge carriers.

In optical generation, charge carriers can have kinetic energies that significantly exceed the average thermal energy of the equilibrium particles. Charge carriers quickly transfer their excess energy to defects in the crystal lattice as a result of propagation. Suppose the energy of the electrons formed by the action of light in the conduction zone is 1 eV .

The electrons are scattered in acoustic phonons with large wavelengths, and at each collision the electrons transmit $2 \cdot 10^{-3} \text{ eV}$ of energy to the grid.

Thus, in order to dissipate energy of a plus of 1 eV , the redundant electrons would have to have about 500 collisions with the acoustic phonons. The average free travel length of electrons is 10^{-6} cm , and their heat velocity at room temperature is approximately 10^7 cm/s , so the mean time between two collisions is $\tau = 10^{-13}, 5 \cdot 10^{-11} \text{ s}$. After that, the excess electrons behave at a crystal lattice temperature. Therefore, the distribution in terms of the energies of the unbalanced and equilibrium charge carriers will be the same. This process is due to the fact that unbalanced extenders dissipate excess energy, as it were, "fall" to the lower edge of the superconductivity zone, and unbalanced spurs "rise" to the upper edge of the stress zone, dissipating excess energy. If, in this case, the concentration of

the unbalanced charge carriers does not differ much from that of the equilibrium, then the transfer of kinetic energy by the unbalanced charge carriers into the crystal does not change the lattice energy as well as the temperature of the crystal, so the concentration of the charge carriers in equilibrium does not change. In this case, the total number of electrons and canopies will be equal to the following,

respectively:

$$n = n_0 + \Delta n; p = p_0 + \Delta p \quad (7.3)$$

Stationary of Non-Balanced Circuits n concentration can be expressed using the formula (4.40):

$$n = n_0 + \Delta n = \sum_{E_C}^{\infty} f_n(E) N(E) dE \quad (7.4)$$

where is the probability of finding an electron at the energy level $f_n = (e^{(E-F_n)/kT} + 1)^{-1}$ E in the conductivity zone. In this case, f_n corresponds to the Fermi-Dirac distribution (4.28), while F_n is the quasi-Fermi level for electrons. (4.42) non-equilibrium electron concentration per expression

$$n = N_C F_{1/2}(\eta_n) \quad (7.5)$$

here

$$\eta_n = \frac{F_n - E_C}{kt} \quad (7.6)$$

The quasi-Fermi level cited for electrons.

Taking into account (4.65), (4.66), and (4.69), the concentration of non-equilibrium electrons in an unseparated semiconductor can be expressed as follows

$$n = N_C e^{\eta_n} = n_0 e^{\eta_n - \eta_i} = n_i e^{\eta_n - \eta_i} \quad (7.7)$$

Similarly, an unbalanced concentration of hives can be expressed as follows:

$$p = p_0 + \Delta p = \sum_{-\infty}^{E_V} f_p(E) N(E) dE = N_V F_{1/2}(-\eta_p - \varepsilon_i) \quad (7.8)$$

where the $f_p = (e^{(E-F_p)/kT} + 1)^{-1}$ probability of finding a hive at the energy level E in the valent zone. In this, f_p corresponds to the Fermi-Dirac distribution (4.28), while f_p is the quasi-Fermi ratio for the spurs.

$$\eta_p = \frac{F_p - E_V}{kt} \quad (7.9)$$

Quasie Fermi brought for hives.

In the absence of sobriety, the concentration of unbalanced hives can be written as follows, taking into account the formulas (4.67) and (4.69):

$$p = N_V e^{-\eta_p - \varepsilon_i} = p_0 e^{\eta_i - \eta_p} = n_i e^{\eta_i - \eta_p} \quad (7.10)$$

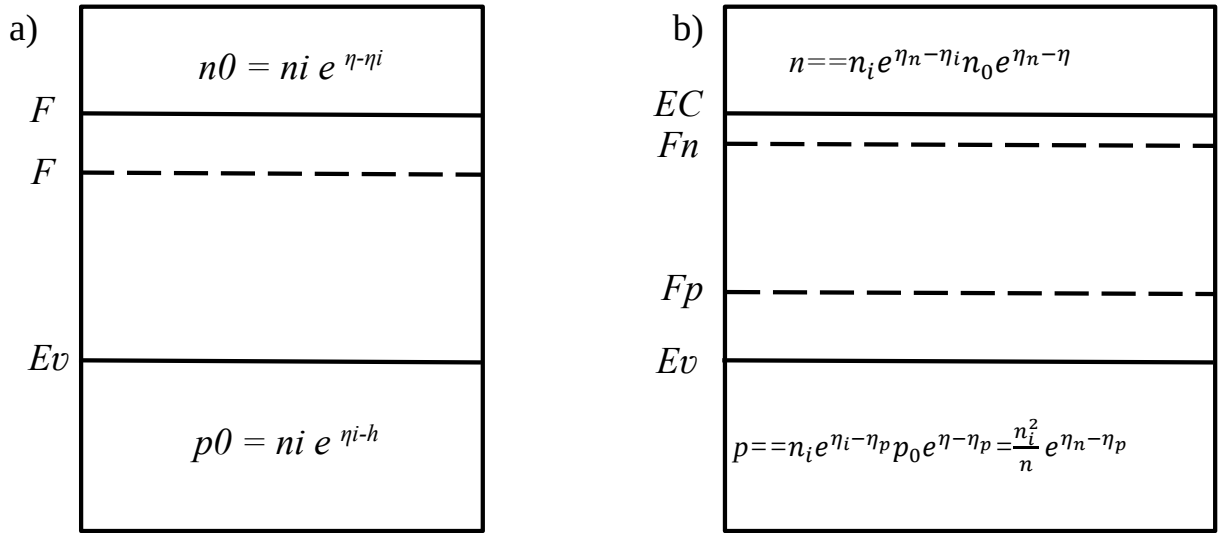


Fig. 7.2. Dividing the Fermi level into two quasi-states: F_n for electrons and F_p for spurs.
(a) state in balance; b) Unbalanced state.

Thus, in an unbalanced state, the Fermi state is divided into two quasi-states, F_n for electrons and F_p for swells, each of which begins to shift towards its own zone (Fig. 7.2).

The multiplication of the concentration of electrons and spurs for the nonequilibrium state is now different from the value at equilibrium:

$$np = n_0 p_0 e^{\eta_n - \eta_p} = n_i^2 e^{\eta_n - \eta_p} \quad (7.11)$$

The distance between the quoted quasi-Fermi levels describes $(\eta_n - \eta_p)$ the deviation of the system from its thermodynamic equilibrium state:

$$np / n_0 p_0 = np / n_i^2 = e^{\eta_n - \eta_p} \quad (7.12)$$

This ratio represents the correlation between the concentration of electrons in the unbalanced state and the fins. In addition, the more different the quasi-Fermi levels of electrons and spurs, the more different the multiplication of the concentration of the unbalanced charge carriers is from the multiplication of the concentration in equilibrium.

7.2. Bipolar optical generation of charge carriers

Consider a semiconductor from which a certain concentration of electrons and spikes is formed by optical generation. We call the bipolar generation of the electron and spur pairs as a result of excitation, the generation of charge carriers. In this case,

the absorption of the quantum of light occurs together with the breaking of the covalent bond, and the number of excess electrons and spurs formed is the same:

$$\Delta n = \Delta p \quad (7.13)$$

In parallel with the generation process, the process of recombination occurs, and the number of electrons and holes excited per unit of time as a result of the breaking of covalent bonds at the stationary state becomes equal to the number of electrons and holes recombined at the same time. Since non-equilibrium charge carriers have become energy-inseparable from equilibrium in a very short period of time, it can be assumed that they have the same recombination coefficient p as equilibrium charge carriers γ .

As soon as the excited light is turned off, the concentration of electrons and the holes by recombination decreases. In this case, the rate of decrease in the number of free electrons and spurs is determined by the difference in the rates of recombination and heat generation

$$-\left(\frac{dn}{dt}\right)_r = -\left(\frac{dp}{dt}\right)_r = \gamma_r np - G_0 \quad (7.14)$$

Taking into account (7.1), (7.2), (7.3) and (7.13), the parity can be rewritten as follows:

$$-\left(\frac{dn}{dt}\right)_f = \gamma_r(np - n_0 p_0) = \gamma_r(n_0 + p_0 + \Delta n)\Delta n \quad (7.15)$$

If the rate of excitation is low, $\Delta n \ll (n_0 + p_0)$ let's take this into account

$$\left(\frac{dn}{dt}\right)_r = -\gamma_r(n_0 + p_0)\Delta n = -\gamma_r(n_0 + p_0)(n - n_0) \quad (7.16)$$

Specify introducing:

$$\tau = \frac{1}{\gamma_r(n_0 + p_0)} \quad (7.17)$$

At the same time, (7.16) takes the following view:

$$\left(\frac{dn}{dt}\right)_r = -\frac{(n - n_0)}{\tau} = -\frac{\Delta n}{\tau} \quad (7.18)$$

here

$$\Delta n = \Delta n(0)e^{-t/\tau} \quad (7.19)$$

in this $\Delta n(0)$ is an excessive concentration of electrons when the excited light is turned off.

Thus, in the case of a low level of bipolar generation, the excess concentration of electrons and holes after the excitation is stopped decreases according to the exponential law, and their number decreases e times over time as a result of recombination. In this case, τ refers to the mean time of existence of the total concentration of excess electrons and spurs, and therefore it is called the lifetime of unbalanced charge carriers. For a private semiconductor, the values of the velocities

of electrons and the number of holes are equal, and the value of τ represents the lifetime of the electron pairs. Depending on the type of semiconductor material and its level of purity, the value of the residence time of unbalanced charge carriers can vary over a very wide range, from 10^{-2} s to 10^{-8} s.

In the case of a high degree of bipolar excitation, we $\Delta n \gg (n_0 + p_0)$ find from (7.15):

$$-\left(\frac{dn}{dt}\right)_r = \gamma_r (\Delta n)^2, \quad (7.20)$$

That is, the rate of recombination Δn will depend on the law of the square to. Breaking down the variables, we write:

$$\frac{dn}{(\Delta n)^2} = -\gamma_r dt, \quad (7.21)$$

here

$$\Delta n = \frac{\Delta n(0)}{1 + \gamma_r \Delta n(0)t} \quad (7.22)$$

It follows from (7.22) that in quadratic recombination, the excess concentration of charge carriers decreases according to the hyperbolic law.

If we enter the instantaneous residence time τ_{on} , in the case of quadratic recombination according to (7.20)

$$\left(\frac{dn}{dt}\right)_r = -\frac{\Delta n}{\tau_{on}} \quad (7.23)$$

and

$$\tau_{on} = \frac{1}{\gamma_r \Delta n} \quad (7.24)$$

That is, τ_{on} depends on the concentration of these excess electrons, and therefore it is a variable magnitude. However, every minute of time τ_{on} has a certain meaning, and a moment of chargers is the time to live. Therefore, according to the high level of injection (7.23), the time of a moment of residence is as follows:

$$\tau_{on} = -\frac{\Delta n}{\frac{d\Delta n}{dt}} \quad (7.25)$$

7.3. Monopolar optical generation of charge carriers. Maxwell Relapse Time

The processes for forming unbalanced charge carriers in input semiconductors can differ significantly from those that occur in privately conductive matter. Consider a donor semiconductor where the illuminated radiation causes electrons to travel from the donor wire to the conduction zone. This process of excitation of charge carriers is called monopolar optical generation. It is

characterized by the formation of an excess concentration of the main charge carriers (being a donor, electrons become the main charge carriers). At the same time, the electrical neutrality of the semiconductor is not disturbed, since the excess charge of free electrons formed by the result of light is compensated by the formed positive charge of the donor crystal ions.

If irequilibrium electrons are formed in a certain region of the sample, then the concentration of electrons in this region increases: $n = n_0 + \Delta n$. The presence of excessive concentration of electrons causes them to pass into the unlumined field of the semiconductor. Therefore, in the unlit part of the sample, where the electrons are diffused, a negative volumetric charge of excess electrons occurs, and in the region where the electrons are diffused, a positive charge accumulates due to donor input ions. you have time $t=0$ When the generation of electrons stops, the electric field E generated by the volumetric charge generates a conductive current that causes the volumetric charge to be eliminated for some time.

The change in the volumetric charge density ρ as a result of a current flow with a density equal to J is known to be subject to the continuity equation of electric charge.

$$\text{div}J = \text{div}(\sigma E) = -\frac{\partial \rho}{\partial t} \quad (7.26)$$

Electric field strength and volumetric charge density are linked by the Poisson equation

$$\text{div}E = \frac{\rho}{\varepsilon_r \varepsilon_0} \quad (7.27)$$

where ε_r is the relative dielectric conductivity of a semiconductor; ε_0 - Electrical constant.

Equations (7.26) and (7.27) show that,

$$\frac{\partial \rho}{\partial t} = \frac{\sigma}{\varepsilon_r \varepsilon_0} \rho \quad (7.28)$$

From here we find the law of change of volumetric charge over time:

$$\rho = \rho_0 e^{-t/\tau_\mu} \quad (7.29)$$

where is the density of the volumetric charge at the moment $\rho_0 - t=0$;

$$\tau_\mu = \frac{\varepsilon_r \varepsilon_0}{\sigma} \quad (7.30)$$

This expression is called dielectric or Maxwell relaxation time.

We estimate the finite value of Maxwell relaxation time. For German

$$\sigma = 1 \frac{1}{\text{Om} \cdot \text{sm}} = 10^2 \frac{1}{\text{Om} \cdot \text{m}} \text{ and } \varepsilon_r = 16, \tau_\mu = 10^{-12} \text{ s}.$$

As can be seen from (7.29), in the case of monopolar optical generation of charge carriers, after the excited light is turned off, the time τ changes according to the exponential law over time. τ_μ A decreasing volumetric charge occurs. In other words, the volumetric charge generated by the excess main charge carriers is the result of a conductive current after the excited light is turned off. T τ_μ disappears after time .

7.4. Recombination Mechanisms

The recombination processes of unbalanced charge carriers in semiconductors can be divided into two main types: interband recombination and recombination by trap.

Interband recombination is accomplished by the transition of an electron from the state in the conductivity band to the empty state of the tensile band, which is equivalent to the destruction of a free electron and a free pit. This process takes place while observing the laws of conservation of energy and quasi-impulse. If the energy and quasi-impulse of an electron in the conductivity zone (before recombination) are E' and p' , and after recombination they are equal to E and p , respectively, then the energy and quasi-impulses are written according to the laws of conservation as follows:

$$E' = E + \Delta E, \quad p' = p + Q \quad (7.31)$$

Q here is the quasi-impulse transmitted by the electron into the crystal lattice and the energy that is released during ΔE -recombination.

Depending on how the excess energy ΔE is distributed, interzone recombination is divided into three types:

- If the energy emitted during recombination is emitted as a quantum of light, the recombination is called radiation or photon recombination;
- If the energy released during recombination is transferred to the lattice, that is, spent to form phonons, such recombination is called non-radiating or phonon;
- If the energy released in the act of recombination is transferred to a third free charge carrier, the recombination is called a shock or Auger recombination.

In a semiconductor crystal there will always be defects, so that their energy layers are located in the forbidden zone. Therefore, as well as inter-zone recombination, a process of recombination through local energy levels can occur. A lattice defect in which it captures an electron from the conductivity zone and a hole from the stress zone, and performs their recombination, is called a recombination trap. In this case, the loss of the circuit and hole pair in conductivity is determined as follows.

The neutral recombination trap pulls the electron out of the conductivity zone, then after a while the electron passes through the stress zone to the stress zone. Depending ΔE on the way the electron is expended when the electron moves to a low-energy state, the recombination through the traps can be photon or phonon.

7.5. Interband Radiation Recombination

During interband radiation recombination, an electron passes through the band of conductivity into the tensile zone, at the same time a quantum of light is released, the energy of which is equal to the width of the bandwidth of the forbidden band:

$$h\nu = E_g \quad (7.32)$$

With the release of photons in thermodynamic equilibrium, the number of recombinable charge carriers in a unit unit of time is equal to the number of charge carriers excited during photon absorption according to the detailed equilibrium principle. In other words, the density of the equilibrium radiation is equal to the R_n^0 number of quanta absorbed in a unit size per second R_y^0 . But the intensity of the radiation recombination process is proportional to the increase in the concentration of electrons and holes in equilibrium:

$$R_y^0 = R_{ch}^0 = \gamma_r n_0 p_0 = \gamma_r n_i^2 \quad (7.33)$$

Here γ_r is the coefficient of interband radiation recombination.

Deviating from the equilibrium state, the unequilibrium is equal to the concentrations of the charge carriers $n = n_0 + \Delta n$ and $p = p_0 + \Delta p$ the equilibrium. They are characterized by the same recombination coefficient as equilibrium, as noted in paragraph 7.2. In this case, the equilibrium of the recombination rate will be maintained in excess of the concentration of unbalanced electrons and spikes, if rotation of the charge carriers does not occur. Therefore, by defining r from (7.33), γ it is possible to write the following:

$$R = \gamma_r np = R_y^0 \frac{np}{n_i^2} \quad (7.34)$$

We determine the residence time of unbalanced charge carriers during interband radiation recombination. For electrons according to (7.18):

$$\tau_{nr} = -\frac{\Delta n}{dn/dt} \quad (7.35)$$

When the external excitation is removed, the rate of change of the concentration of free electrons is determined by the difference of the recombination intensity P and the equilibrium generation R_y^0 . At the same $R_y^0 = R_n^0$ time, without

$$-\frac{dn}{dt} = R - R_n^0 \quad (7.36)$$

Given this, (7.35) is written as follows

$$\tau_{nr} = \frac{\Delta n}{R - R_n^0} = \frac{\Delta n}{\Delta R} \quad (7.37)$$

Here ΔR is the change in the rate of recombination when the system deviates from the equilibrium state.

A similar equation can be written for holes:

$$\tau_{pr} = \frac{\Delta p}{R - R_n^0} = \frac{\Delta p}{\Delta R} \quad (7.38)$$

We define the amount P using the expressions (7.34) and (7.33):

$$\Delta R = R - R_n^0 = R_n^0(np - n_i^2)/n_i^2 \quad (7.39)$$

By putting the values n and p in (7.39) and replacing n_i^2 it by $n_0 p_0$, we get:

$$\Delta R = R_n^0 \frac{n_0 \Delta p + p_0 \Delta n + \Delta n \Delta p}{n_0 p_0} \quad (7.40)$$

Now (7.37) takes the following view:

$$\tau_{nr} = \Delta n \frac{n_0 p_0}{R_n^0 (n_0 \Delta p + p_0 \Delta n + \Delta n \Delta p)} \quad (7.41)$$

If $\Delta n = \Delta p$, (7.41) is simplified:

$$\tau_{nr} = \tau_{pr} = \frac{n_0 p_0}{R_n^0 (n_0 + p_0 + \Delta n)} \quad (7.42)$$

Let us analyze the resulting expression. If the degree of excitation is small $\Delta n \ll (n_0 + p_0)$ Taking into account (7.33):

$$\tau_{nr} = \tau_{pr} = \frac{n_i^2}{R_n^0 (n_0 + p_0)} = \frac{1}{\gamma_r (n_0 + p_0)} \quad (7.43)$$

For Private Semiconductor $n_0 = p_0 = n_i$,

$$\tau_{ir} = \frac{n_i}{2R_n^0} = \frac{1}{2\gamma_r n_i} = \frac{e^{E_g/2kT}}{2\gamma_r (N_C N_V)^{1/2}} \quad (7.44)$$

n-for different material ($n_0 \gg p_0$)

$$\tau_r^n = \frac{p_0}{R_n^0} = 2 \frac{n_i}{n_0} \tau_{ir} = \frac{1}{\gamma_r n_0} \quad (7.45)$$

and for the hollow semiconductor ($p_0 \gg n_0$)

$$\tau_r^n = \frac{n_0}{R_n^0} = 2 \frac{n_i}{p_0} \tau_{ir} = \frac{1}{\gamma_r p_0} \quad (7.46)$$

Analysis of the (7.44)-(7.46) formulas shows that the residence time of non-equilibrium charge carriers in a private semiconductor during interband radiation recombination is greater than the temperature and the smaller the bandwidth (see Table 7.1). The residence time of non-equilibrium charge carriers in an input

semiconductor is much smaller than that of a private semiconductor, and it decreases as the rate of discharge and temperature increase τ_{ir} .

The dependence of residence time on the level of lesion, i.e., the concentration of charge carriers at the small excitation level, constant temperature, for interband radiation recombination is shown in Fig. 7.3. In this figure, the logarithmic scale of the charge carrier concentration can be considered to be the same as the linear scale of the Fermi surface state, since $n = N_c e^{-(E_c - F)/kT}$, moreover, its midpoint corresponds to the value of the Fermi level for a private semiconductor.

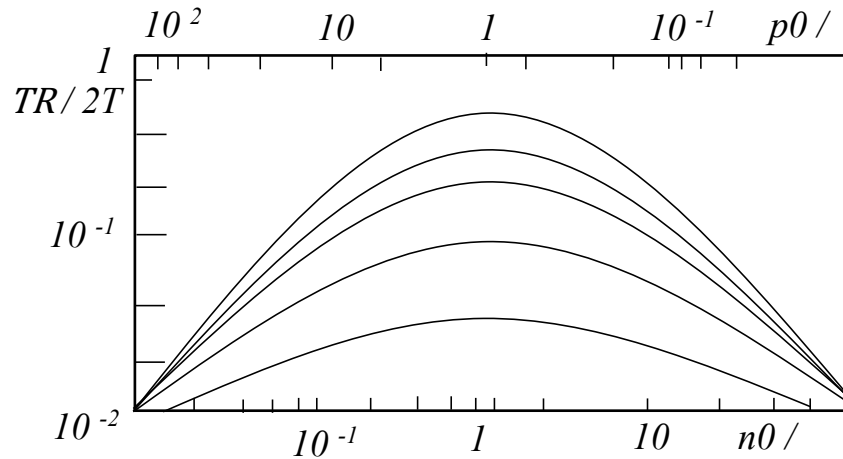


Fig. 7.3. Residence time for interband radiation recombination depends on concentration of charge carriers at a given temperature at low excitation rate. The ratio for the maximum is on the upper curve $\Delta n / n_i = 0$, for the other subsequent curves it is 1, 3, 10, 30, respectively

From the above dependencies, it follows that with the increase in the excitation rate, the residence time in the private semiconductor decreases dramatically, and the residence time in the input semiconductor changes relatively slowly.

R_n^0 Using Einstein's theory of radiation in equilibrium. The $P + dp$ number of quantum states in the p and impulse range for photons is:

$$P^2 dP / h^3 \quad (7.47)$$

Is equal to.

To determine the number of cells occupied by photons, this value must be multiplied by the Bose-Einstein distribution function:

$$f_f = \frac{1}{e^{h\nu/kt} - 1} \quad (7.48)$$

The Ω total number of photons with a P impulse emitting at an angle d is as follows:

$$dn_f = \frac{2}{h^3} f_f P^2 dP d\Omega \quad (7.49)$$

where the coefficient of 2 takes into account the different polarity of the light. The photon impulse is related to energy as follows

$$(E = h\nu);$$

$$P = E/\nu_f = h\nu/\nu_f = nh\nu/c \quad (7.50)$$

in this sweat $\nu_f = c/n$ - the speed of a photon in a given medium; c -The speed of light in a vacuum; n - The refractive index of the medium.

(7.50) and:

$$dP = \frac{nh}{c} d\nu \quad (7.51)$$

By putting this expression in (7.49), we get the following expression, which is integrated through all angles from 0 to 4π :

$$dN = \int_0^{4\pi} \frac{2}{h^3} f_f P^2 dP d\Omega = \frac{8\pi}{h^3} f_f P^2 dP = \frac{8\pi}{c^3} \frac{n^3 \nu^2 d\nu}{e^{h\nu/kT} - 1} \quad (7.52)$$

If ν the probability of absorption of a frequency photon $g(\nu)$ We get the following expression for absorption intensity equal to the intensity of radiation under thermodynamic equilibrium:

$$R_n^0 = R_y^0 = \int_0^\infty g(\nu) dN = \frac{8\pi}{c^3} \int_0^\infty \frac{n^3 g(\nu) \nu^2 d\nu}{e^{h\nu/kT} - 1} \quad (7.53)$$

Photon's Lifetime τ_f gives the inverse of the probability of its absorption. free-treadmill length of magnitude $\tau_f l_f$ and speed ν_f can be found when known. At the same time, the length of the free-running path l_f absorption coefficient $\alpha(\nu)$ Not just for the Tesla, yes. $l_f = 1/\alpha(\nu)$. As a result, we get:

$$\tau_f = 1/g(\nu) = l_f / \nu_f = n/c\alpha(\nu) \quad (7.54)$$

$g(\nu)$ By placing the value of (7.54) and (7.53) we find:

$$R_n^0 = \gamma_r n_i^2 = \frac{8\pi}{c^2} \int_0^\infty \frac{n^2 \alpha(\nu) \nu^2 d\nu}{e^{h\nu/kT} - 1} \quad (7.55)$$

Table 7.1 shows the values and lifetime values $n_i^2 \tau_{ir} R_n^0$ for the radiated conduction zone-stress band transitions in some semiconductors. Here τ_{eks} is the experimentally measured value. The data in Table 7.1 show that in the case of inter-band radiation recombination, the residence time of non-equilibrium charge carriers decreases with the reduction of the bandwidth of the bandwidth. The calculated τ_{ir} values are close to the result obtained only in experiments from semiconductors with a small bandwidth limit. It is observed that the data calculated for substances with a large forbidden zone width differ significantly with the experimentally measured results. This is due to the high probability of a radiation-free transition in these semiconductors.

Table 7.1. Radiation recombination in some semiconductors at 300 K temperatures

Semiconductor	Eg, eV	Ni, ^{SM-3}	R_n^0 , $\text{sm}^{-3} \cdot \text{s}^{-1}$	tir, mks	Teeks, MKS
Has					
GaAs	1,43	$4,3 \cdot 10^{12}$	$2,4 \cdot 10^{14}$	9000	0.37
InAs	0,31	$1,6 \cdot 10^{15}$	$5,4 \cdot 10^{19}$	15	0.24
InSb	0,18	$2,0 \cdot 10^{15}$	$1,6 \cdot 10^{22}$	0.62	0.12
PbS	0,41	$7,1 \cdot 10^{14}$	$2,4 \cdot 10^{19}$	15	0.21
PbTe	0,32	$4,0 \cdot 10^{15}$	$8,3 \cdot 10^{20}$	2.4	0.19
PbSe	0,29	$6,2 \cdot 10^{15}$	$1,5 \cdot 10^{21}$	2	0.25

According to the formulas (7.45) and (7.46), the residence time of unbalanced charge carriers decreases with an increase in the level of legging. For example, for strongly alloyed germanium ($n_o = 10^{19} \text{sm}^{-3}$), this is:

$$\tau_r^n = 2n_i \tau_{ir} / n_0 \approx 2 \cdot 2,4 \cdot 10^{13} \cdot 0.61 / 10^{19} \approx 3 \cdot 10^{-6} \text{s} = 3 \text{mks}$$

As seen above, as the degree of lesion of the material increases, the possibility of inter-band radiation recombination must increase. However, with an increase in the concentration of inputs, many defects appear and recombination through traps begins to dominate. Therefore, the purest semiconductor materials need to be used to study the recombination processes of interband radiation.

7.6. Interzonal Shock Recombination

If a collision occurs simultaneously between two free electrons and one pit, or between two hives and one free electron, the recombination of two of them can occur by transferring energy to the third charge carrier, while the third charge passes to a high-energy state in the corresponding zone and becomes "hot".

The "hot" carrier transfers its energy to the phonons as a result of collisions and comes into a state of equilibrium with the grid. This recombination process is the inverse of shock ionization, in which the transition from the free charge carrier in the zone to the low energy spike as a result of energy release occurs when an electron hive pair is formed.

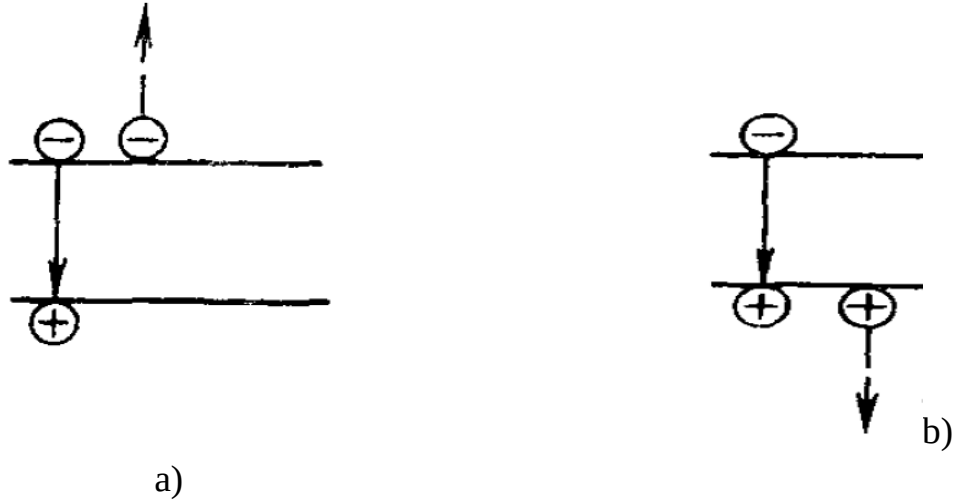


Fig. 7.4. Scheme of interband pulse recombination processes

In Fig. 7.4, the electron (7.4, Fig. a) or the electron (7.4, Fig. A) involved as a third charge carrier *b* rasm) The shock recombination processes of On are schematically described.

The probability of a collision of an electron pair of an electron with a free electron is proportional to n^2p , and the probability of a collision with a pair of electrons is proportional to n^2p . Thus, the decrease in the concentration of free electrons or spikes as a result of interband shock recombination is determined by the following expression.

$$-\frac{dn}{dt} = -\frac{dp}{dt} = \eta^{(n)} n^2 p + \eta^{(p)} p^2 n - \eta^{(n)} n_0^2 p_0 - \eta^{(p)} p_0^2 n_0 \quad (7.56)$$

here $\eta^{(n)}, \eta^{(p)}$ are the coefficients of the shock recombination of the electron or the spur that acts as the third charge carrier. The last two figures in (7.56) take into account the intensity of the generation in thermodynamic equilibrium. It is important to note that shock recombination coefficients differ measurably from the corresponding coefficients of other types of recombination.

By putting the values n and p at (7.56), equalizing $\Delta n = \Delta p$, and putting the equation states in Δn degrees, we get:

$$-\frac{d}{dt} \Delta n = -\frac{d}{dt} \Delta p = \Delta n (\eta^{(n)} n_0^2 + 2\eta^{(n)} n_0 p_0 + 2\eta^{(p)} n_0 p_0 + \eta^{(p)} p_0^2) + \Delta n^2 (\eta^{(n)} p_0 + 2\eta^{(n)} n_0 + 2\eta^{(p)} p_0 + \eta^{(p)} n_0) + \Delta n^3 (\eta^{(n)} + \eta^{(p)}) \quad (7.57)$$

Assuming that the injection rate is small $\Delta n \ll (n_0 + p_0)$, we find an expression for the residence time of unbalanced charge carriers in the case of interband shock recombination in accordance with (7.18). To do this, we throw out the hads from (7.57) that contain in the second and third degrees Δn , and we get the following:

$$\tau_{nA} = \tau_{pA} = -\frac{\Delta n}{\frac{d}{dt}\Delta n} = \frac{1}{\eta^{(n)}n_0^2 + 2\eta^{(n)}n_i^2 + 2\eta^{(p)}n_i^2 + \eta^{(p)}p_0^2} \quad (7.58)$$

We will consider specific cases. 1. Private semiconductor ($n_0 = p_0 = p_i$) To:

$$\tau_{iA} = \frac{1}{3(\eta^{(n)} + \eta^{(p)})n_i^2} = \frac{1}{3(\eta^{(n)} + \eta^{(p)})N_C N_V} e^{E_g / kT} \quad (7.59)$$

It can be seen from the expression that the residence time of the unbalanced charge carriers in a private semiconductor decreases with temperature increase and the bandwidth of the bandwidth decreases in the interband shock recombination.

2. For the donor semiconductor ($n_0 \gg p_0; n_0 \gg n_i$)

$$\tau_A^n = \frac{1}{\eta^{(n)}n_0^2} \quad (7.6)$$

3. Actuator for semiconductor ($p_0 \gg n_0; p_0 \gg n_i$)

$$\tau_A^p = \frac{1}{\eta^{(p)}p_0^2} \quad (7.61)$$

From the obtained expressions, it follows that the residence time of non-equilibrium charge carriers is maximum for a private semiconductor in interband charge recombination. In this type of recombination in an input semiconductor, the residence time of the non-equilibrium charge carriers is very strongly dependent on the concentration of the main charge carriers. With an increase in the concentration of primary charge carriers, the number of collisions of non-equilibrium electronic hve pairs with primary charge carriers increases and the residence time decreases. τ_A / τ_i The dependence of n_i on the state of the Fermi level is shown in Fig. 7.5.

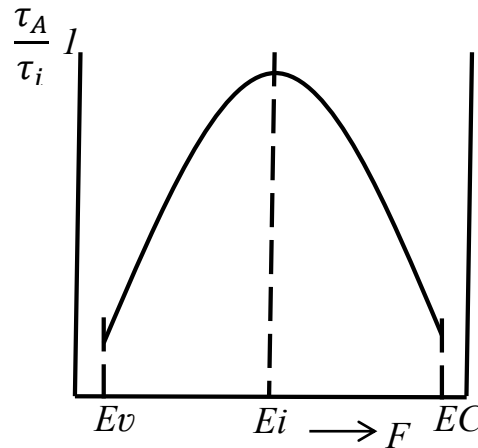


Fig. 7.5. Dependence of residence time on Fermi level state at the time of interzonal shock recombination

Experimental studies have shown that interzone forgiveness recombination is extremely rare. It is noticeable in semiconductors (e.g. InSb) where the bandgap width is small at very high temperatures.

At a high excitation level, $\Delta n \gg (n_0 + p_0)$ and $\Delta p \gg (n_0 + p_0)$ that can be considered. For this case, the first two had in (7.57), i.e., Δn and Δn^2 those containing, can be discarded. As a result, we get:

$$-\frac{d}{dt}\Delta n = -\frac{d}{dt}\Delta p = \Delta n^3(\eta^{(n)} + \eta^{(p)}) \quad (7.62)$$

In this case, the time of instantaneous (temporary) residence

$$\tau_{oniy} = \frac{\Delta n}{\frac{d}{dt}\Delta n} = \frac{1}{(\eta^{(n)} + \eta^{(p)})\Delta n^2} \quad (7.63)$$

From (7.63), it follows that the instantaneous residence time in an interzonal charge recombination at a high excitation level is strongly dependent on the concentration of excess charge carriers.

7.7. Recombination of Charge Carriers by Traps

True semiconductor materials typically contain several types of inputs, each of which can generate one or more energy layers in the semiconductor prohibited area.

Lattice defects, which are usually neutral in a thermodynamic equilibrium state and are able to grab and release actuating moving charge carriers, are called traps or capture centers. We will limit ourselves to considering one of the simplest types of traps that form a single energy line in a semiconductor.

Let the concentration of the traps be equal to N_t and their energy level E_t be located near the middle of the prohibited zone. The process of recombination through these traps is shown in Fig. 7.6.

In a donor-type semiconductor with a large number of free electrons, the recombination of an electron and a hole is as follows. A neutral trap acquires a negative charge by capturing the electron conductivity zone (transition 1), and then passes from the electron trap surface to the tensile zone, which is equivalent to catching a snare with a negatively charged trap (transition 2). With a decrease in the concentration of free holes, the probability and electron emission of the second process, i.e. the reverse circuit of the electron from the trap to the permeability zone (transition 3), decreases.

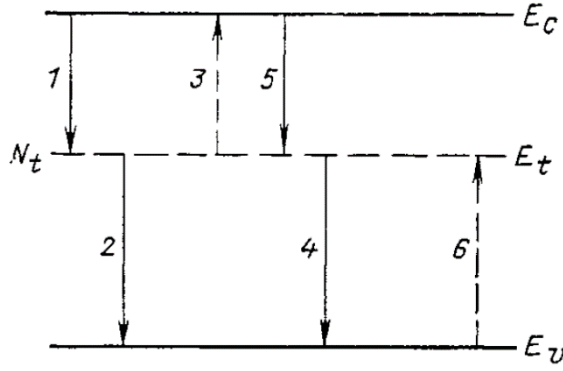


Fig. 7.6. Electron crossovers associated with recombination through the trap

If there are a large number of free swells in the material, they are intensely captured by traps, that is, electrons pass through the trap into the stress zone (transition 4), and then capture of the electron in the conduction zone occurs (transition 5). At a small concentration of conduction electrons, the process of hive release can occur, i.e., it is possible to return the hive back to the tensile zone, (transition 6).

We will move on to a quantitative description of the processes considered. If $f_t = f(E)$ we define the probability of filling the trap with an electron, then $1 - f_t$ there will be the probability that there is no electron in the trap, that is, the probability of electron freedom. In this case, the concentration of electrons in the conduction zone is proportional to n and the $N_t(1 - f_t)$ intensity of electron capture by the handles is proportional to the number of empty spaces at the handle level:

$$R_{ut}^n = \gamma_n n N_t (1 - f_t) \quad (7.64)$$

Here γ_n is the coefficient of traps catching unbalanced electrons.

The number of electrons returning to the permeability zone as a result of emission from R_{qay}^n traps is proportional to the concentration of electrons in the traps $N_t f_t$.

$$R_{qay}^n = \beta_n N_t f_t \quad (7.65)$$

Here β_n is the coefficient of ionization of electrons from traps.

The change in the concentration of non-equilibrium electrons in the permeability zone R_{ut}^n is determined by the difference between R_{qay}^n the intensity of the compression processes and the intensity of release:

$$-\frac{dn}{dt} = R_{ut}^n - R_{qay}^n = \gamma_n n N_t (1 - f_t) - \beta_n N_t f_t \quad (7.66)$$

γ_n and β_n To determine the relationship between the coefficients, we look at the thermodynamic equilibrium condition. The probability of electron filling the trap level in the thermodynamic equilibrium state:

$$f_{0t} = f_0(E_t) = \frac{1}{e^{(E_t - F)/kT} + 1}$$

If spinning does not occur in the unbalanced state, then most of the time that unequilibrium charge carriers live is indistinguishable from those in equilibrium. Therefore, we can assume that charge carriers in equilibrium are also γ_n characterized by both retention and ionization β_n coefficients.

In the thermodynamic equilibrium state, for (7.66) $\frac{dn}{dt} = 0$ Based on We find:

$$\beta_n = \gamma_n n_0 \left(\frac{1}{f_{0t}} - 1 \right) = \gamma_n N_c e^{-(E_c - F)/kT} e^{(E_t - F)/kT} = \gamma_n N_c e^{-(E_c - E_t)/kT} = \gamma_n n_1 \quad (7.67)$$

here

$$n_1 = N_c e^{-(E_c - E_t)/kT} \quad (7.68)$$

The expression gives the equilibrium concentration of electrons in the conductivity zone when the Fermi level corresponds to the trap level.

Using (7.67), we pass equation (7.66) to:

$$-\frac{dn}{dt} = \gamma_n N_t [n(1 - f_t) - n_1 f_t] \quad (7.69)$$

This expression determines the rate of change of the non-equilibrium electron concentration in the conductivity zone as a result of the processes of entrapping electrons and releasing them from the traps into the conductivity zone.

The rate of change in the concentration of unbalanced canines in the stress zone is written as follows:

$$-\frac{dp}{dt} = \gamma_p p N_t f_t - \beta_p N_t (1 - f_t) \quad (7.70)$$

Here, the first determines the number of hives caught from the traps, and the second determines the number of hives that return from the traps to the tension zone. γ_p and β_p The magnitudes represent the coefficients of capture and ionization, respectively, for hives.

From the thermodynamic equilibrium we $\frac{dp}{dt} = 0$ find:

$$\beta_p = \gamma_p p_0 \frac{f_{0t}}{1 - f_{0t}} = \gamma_p N_v e^{(E_v - F)/kT} e^{-(E_t - F)/kT} = \gamma_p N_v e^{(E_v - E_t)/kT} = \gamma_p p_1 \quad (7.71)$$

here

$$p_1 = N_v e^{(E_v - E_t)/kT} \quad (7.72)$$

Expression The equilibrium concentration of hives in the tensile zone when the Fermi level coincides with the trap level.

If we replace (7.71) with (7.70), we get the following expression:

$$-\frac{dp}{dt} = \gamma_p N_t [pf_t - p_1(1 - f_t)] \quad (7.73)$$

This expression defines the rate at which the concentration of unbalanced hives in the stress zone changes as a result of the processes of trap hives catching and releasing them from the traps into the stress zone.

Using (7.69) and (7.73), we find the variation in the number of electrons located in the traps. It is determined by changes in the concentration of unbalanced electrons in the conductivity zone and cavities in the stress zone:

$$\frac{dn}{dt} - \frac{dp}{dt} = \gamma_p N_t [pf_t - p_1(1 - f_t)] - \gamma_n N_t [n(1 - f_t) - n_1 f_t] \quad (7.74)$$

N_d^+ We write a local electrical neutrality scenario for a single-gender semiconductor containing a certain amount of donor input ions. The electric neutrality condition for the equilibrium state is as follows:

$$n_0 + N_t f_{0t} = p_0 + N_d^+ \quad (7.75)$$

We estimate and assume that the number of donor input ions does not change even in the imbalanced state. Without it:

$$n + N_t f_t = p + N_d^+ \quad (7.76)$$

From (7.76) and (7.75) we get the relation for the excess concentration of charge carriers:

$$n - n_0 + N_t (f_t - f_{0t}) = p - p_0$$

or

$$\Delta n + N_t (f_t - f_{0t}) = \Delta p \quad (7.77)$$

If the concentration of traps in the semiconductor is small in relation to the density of the excess charge carriers, that is, $N_t \ll \Delta n, \Delta p$, then Number of electrons located in the traps $N_t \Delta f_t$ ni It is possible to ignore. Based on this condition, the excess concentration of electrons is equal to the excess concentration of the sleds, i.e.

$$\Delta n = \Delta p \quad (7.78)$$

For this case $\frac{dn}{dt} = \frac{dp}{dt}$, and from (7.74), by zeroing it to zero, we find the distribution function for the electrons located in the traps:

$$f_t = \frac{\gamma_n n + \gamma_p p_1}{\gamma_n (n + n_1) + \gamma_p (p + p_1)} \quad (7.79)$$

By putting the found f_t value at (7.69), we get:

$$-\frac{dn}{dt} = \frac{\gamma_n \gamma_p N_t (np - n_1 p_1)}{\gamma_n (n + n_1) + \gamma_p (p + p_1)} \quad (7.80)$$

Based on this ratio, it is possible to determine the residence time of unbalanced charge carriers during recombination by traps.

This

$$n_1 p_1 = n_i^2 = n_0 p_0 \quad (7.81)$$

By taking into account, the result is the following:

$$\tau = -\frac{\Delta n}{\frac{dn}{dt}} = \frac{\gamma_n(n_0 + n_1 + \Delta n) + \gamma_p(p_0 + p_1 + \Delta p)}{\gamma_n \gamma_p N_t (n_0 + p_0 + \Delta n)} = \frac{1}{\gamma_p N_t} \frac{n_0 + n_1 + \Delta n}{n_0 + p_0 + \Delta n} + \frac{1}{\gamma_n N_t} \frac{p_0 + p_1 + \Delta p}{n_0 + p_0 + \Delta n} \quad (7.82)$$

We include the following definitions:

$$\tau_{p0} = \frac{1}{\gamma_p N_t}; \quad \tau_{n0} = \frac{1}{\gamma_n N_t} \quad (7.83)$$

Given (7.83), the formula (7.82) can be written as follows

$$\tau = \tau_{p0} \frac{n_0 + n_1 + \Delta n}{n_0 + p_0 + \Delta n} + \tau_{n0} \frac{p_0 + p_1 + \Delta p}{n_0 + p_0 + \Delta n} \quad (7.84)$$

If there is a small excitation rate, that is, if the $\Delta n \ll (n_0 + p_0)$ condition is met, the residence time of the unbalanced charge carrier depends on the state of the Fermi level at constant temperature, depending on the degree of oxidation of the material, it will be equal to:

$$\tau = \tau_{p0} \frac{n_0 + n_1}{n_0 + p_0} + \tau_{n0} \frac{p_0 + p_1}{n_0 + p_0} \quad (7.85)$$

So, At the small excitation level, the residence time of the electron hole pairs does not depend on the concentration of overcharge carriers, but is determined by the values of the concentration of electrons and holes in equilibrium, and by the state of the recombination trap energy level, which determines the values of n_1 and p_1 .

Based on the formula (7.85), four main areas can be distinguished from the dependence of $\ln \tau$ on the Fermi level state (Fig. 7.7).

1 - Sphere. *N-type* strong alloyed semiconductor. The Fermi level is located at the bottom of the permeability zone, but above the energy level of the recombination trap, i.e., $E_t < F < E_C$. In this case, the following ratio of the concentration of charge carriers applies:

$$n_0 \gg p_0; n_0 \gg n_1; n_0 \gg p_1;$$

Given these disparities, we get the following from (7.85):

$$\tau = \tau_{p0} = \frac{1}{\gamma_p N_t} \quad (7.86)$$

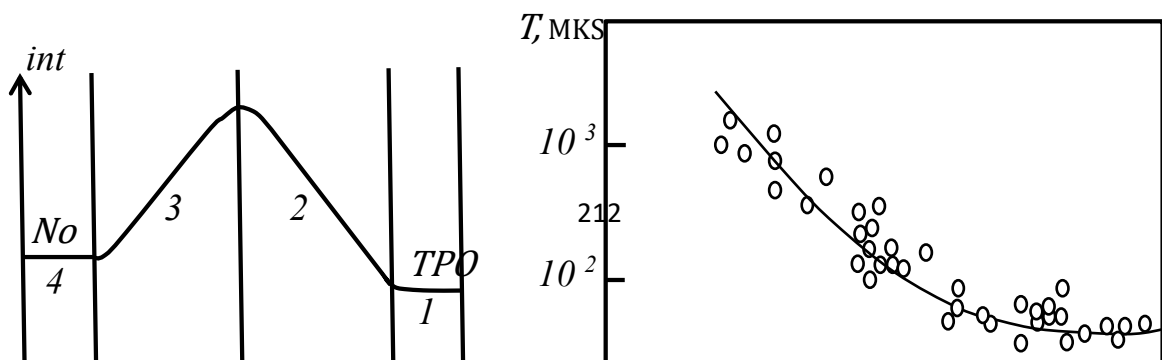


Fig. 7.7. Dependence of the residence time of non-equilibrium charge carriers on the state of the Fermi level

Fig. 7.8. The dependence of the survival time of hives in an electron germanium depends on the concentration of electrons in equilibrium

Thus, in a strongly propelled n-type semiconductor, the lifetime of non-equilibrium electrons and hives is unchanged and equal to $\tau = \tau_{p0}$, i.e., these non-fundamental equilibrium charge carriers, are determined by the lifetime of the hives. In this case, the speed of the recombination process is determined by the number of actions of the traps occupied by the spurs, which are occupied by the spurs. Indeed, experience shows (Fig. 7.8) that the lifespan of hives in a strongly alloyed electron is constant.

(7.86) We analyze the expression. This equation represents the above fact, for example, that when each cavity in an n-type material moves through the stress zone, it is likely to fall into a negatively charged recombination trap. To quantify this process, we use σ_p from an efficient cross-section of hive capture. According to (5.48), the probability of catching an unbalanced charge carrier with a single center is proportional to the effective cross-section of catching the snare with a trap and v_0 to the heat velocity of the charge carriers, hence the following expression:

$$\gamma_p = \sigma_p v_p \quad (7.87)$$

The lifetime of a hive in an electronic semiconductor is as follows:

$$\tau = \tau_{p0} = \frac{1}{\gamma_p N_t} = \frac{1}{\sigma_p N_t v_0} \quad (7.88)$$

This is The recombination traps decrease with increasing sample temperature, resulting in increased concentration and an increase in v_0 .

Area 2. An n-type weakly alloyed semiconductor in which the Fermi level lies in an area below the energy level of the traps but above the middle of the forbidden zone, i.e., $E_i < F < E_t$. In this case, the following relationship is appropriate for the concentration of charge carriers:

$$n_0 \gg p_0; n_0 \gg p_1; n_0 \ll n_1.$$

Accordingly, the expression for τ is:

$$\tau \approx \tau_{p0} \frac{n_1}{n_0} = \tau_{p0} e^{(E_t - F)/kT} \quad (7.89)$$

It follows from (7.89) that when the Fermi level decreases, the residence time of non-equilibrium charge carriers increases according to the exponential law τ . However, the lower the Fermi level, the lower the rate at which the traps are charged with electrons, so the likelihood of a trap catching the hive decreases, which increases the hives' residence time. This is confirmed by the data in Fig. 7.7.

Sphere 3. In a weakly alloyed p-type semiconductor, the state of the Fermi level $(E_v - E_t) < F < E_i$ is determined by inequality, the following relation is appropriate for the concentration of charge carriers:

$$p_0 \gg n_0; p_0 \gg p_1; n_1 \gg p_0.$$

Based on these, we get the following formula for the length of residence from (7.85):

$$\tau \approx \tau_{p0} \frac{n_1}{p_0} = \tau_{p0} \frac{N_C}{N_V} e^{-(E_C + E_V - E_t - F)/kT} \quad (7.90)$$

Ko'rinib turibdiki, When the Fermi level for the accelerator semiconductor falls below the middle of the exclusion zone as a result of logging, the residence time of the non-equilibrium charge carriers decreases exponentially. This is because p In a weak alloyed semiconductor of the -type, almost all of the traps are in the free state and pull electrons out of the conductivity zone. With an increase in the number of permeability hives, the possibility of recombining them with trapped electrons increases. In this case, the intensity of the electron release process from the traps decreases and the residence time of unbalanced charge carriers decreases (Fig. 7.7).

Area 4. A p-type strongly alloyed semiconductor, the state of the Fermi level is determined by the $E_v < F < (E_v - E_t)$ inequality. In this case $p_0 \gg n_0; p_0 \gg p_1; p_0 \gg n_1$

Appropriately, from (7.85) the following is obtained:

$$\tau = \tau_{n0} = \frac{1}{\gamma_n N_t} = \frac{1}{\sigma_n N_t \nu_0} \quad (7.91)$$

Here σ_n is an efficient section of electron capture.

Thus, in a p-type strongly alloyed semiconductor, the lifetime of an electron pair is constant, equal to $\tau = \tau_{n0}$, determined by the residence time of the electron-non-basic charge carriers in equilibrium, and is independent of the state of the Fermi level. In such a semiconductor, all the traps are in the electrons, and each electron captured by the trap is immediately combined with the hollow, because the number of pitfalls in the stress band is very large. The process of returning electrons from traps to the permeability zone is almost non-existent and has no effect on survival.

Now consider the high-level excitation case of excess charge carriers, wherein $\Delta n \gg n_0; \Delta p \gg p_0$.

So,

$$\Delta n \gg n_1; \Delta p \gg p_1$$

will be appropriate. Taking this into consideration, we get the following statement from (7.84):

$$\tau_\infty \approx \tau_{p0} + \tau_{n0} = \frac{\gamma_n + \gamma_p}{\gamma_n \gamma_p N_t} \quad (7.92)$$

From this expression you can see that, Unbalanced electrons at large excitation levels and hives live-timet τ_{∞} When recombined through traps, it does not depend on the concentration of electrons and hives, but is determined only by the number and characteristics of the traps.

7.8. Temperature Dependence of Charge Carrier Residence Time During Recombination Through Traps

Let us consider how the residence time of non-equilibrium charge carriers changes when the research temperature changes during recombination through traps in the semiconductor of a given structure. Let us do this analysis by the example of an electron semiconductor. Let's say that in such an inlet semiconductor, when the temperature rises, as shown in Fig. 7.9, three changes in the concentration of charge carriers are observed, i.e., the area of inlet conductivity (Area A), The total ionization area of the inlet, characterized by the constant concentration of charge carriers (Area B), and private conductivity (Area V).

Temperature direction 7.9, b lies between the Fermi stage in Area A, shown in Fig. A, the level of the donor clamp and the lower part of the permeability zone. As temperatures increase, the Fermi level drops to the donor entry level. Because the Fermi level for all temperatures in the A area is higher than the energy level of the recombination trap, the concentration of electrons

$$n_0 \gg n_1; n_0 \gg p_0; n_0 \gg p_1$$

Subject to inequalities. Consequently, in accordance with (7.86) in this area, the residence time of unbalanced hives is constant and *equal to* τ_{p0} (7.9-v, Field 1 in Fig.1, which shows that the residence time is temperature-dependent).

As the temperature increases further during the transition to area B, the Fermi level decreases in the area of $E_d < F < E_t$, i.e., the complete ionization of the donor input, and reaches the level of the recombination trap.

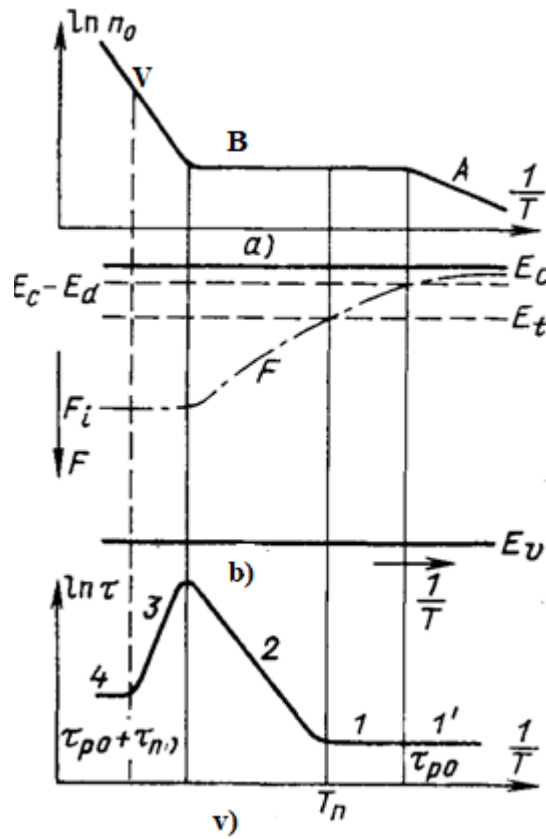


Fig. 7.9. Temperature dependence of concentration (a), Fermi level (b), and charge carrier residence time (v) in equilibrium in the donor semiconductor.

In this field, the concentration of electron n_0 is constant, and the $n_1 = N_c e^{-(E_C - E_t)/kT}$ concentration increases exponentially with temperature, but it still has a concentration at equilibrium that is significantly lower than n_0 . In addition, the concentration of n_0 hives is even greater than p_0 and p_1 . According to the expression in this area (7.86), the residence time is also constant and equal to τ_{p0} (Fig. 7.9, Fig. 1).

With a further increase in temperature, the Fermi level decreases and becomes lower than the energy level of the recombination trap (Sphere 2 in Figure 7.9-v). Here, as before, $n_0 = \text{const}$, and n_1 increases considerably, so that from (7.90) we can write the following expression:

$$\tau \approx \tau_{p0} \frac{n_1}{p_0} = \tau_{p0} \frac{N_c}{n_0} e^{-(E_C - E_t)/kT} = \frac{\tau_{p0}}{n_0} 2 \left(\frac{2\pi m_n^* k}{h^2} \right)^{3/2} T^{3/2} e^{-(E_C - E_t)/kT} \quad (7.93)$$

That is, the residence time of non-equilibrium charge carriers increases with increasing temperature. At the same time, it increases to the temperature at which private permeability occurs (7.9, Fig. v, Sphere 2). In this temperature region, an increase in τ occurs, with an increase in temperature, an intense emission of electrons occurs in the trap. In this case, the electron filling of the traps decreases, and

therefore the life span of the hives, non-primary carriers, which represent the measured time τ , increases.

As temperatures rise in the private conductivity region (in area b), although the state of the Fermi level remains almost constant (Figure 7.9, b), the concentration of charge carriers in the zones increases sharply, so that the filling of traps with charge carriers increases and the residence time of unbalanced electrons and spurs decreases. For this area, the following statement is appropriate:

$$n_0 \approx p_0 \approx n_i = (N_C N_V)^{1/2} e^{-E_g / 2kT}. \quad (7.94)$$

In accordance with the above (7.85) it is written as follows:

$$\begin{aligned} \tau &= \tau_{p0} \frac{n_i + n_l}{2n_l} + \tau_{n0} \frac{n_i + p_l}{2n_l} = \frac{\tau_{p0}}{2} \left[1 + \left(\frac{N_C}{N_V} \right)^{1/2} e^{((E_g / 2) + E_t - E_C) / kT} \right] + \frac{\tau_{n0}}{2} \left[1 + \left(\frac{N_V}{N_C} \right)^{1/2} e^{((E_g / 2) - E_t + E_V) / kT} \right] = \\ &= \frac{\tau_{p0}}{2} \left[1 + \left(\frac{N_C}{N_V} \right)^{1/2} e^{-(((E_C + E_V) / 2) - E_t) / kT} \right] + \frac{\tau_{n0}}{2} \left[1 + \left(\frac{N_V}{N_C} \right)^{1/2} e^{(((E_C + E_V) / 2) - E_t) / kT} \right] \end{aligned} \quad (7.95)$$

If the level of recombination traps is located far enough from the middle of the forbidden zone, then at the very beginning of the private permeability $[(E_C + E_V) / 2 - E_t] \gg kT$ it is possible to isolate a place where unevenness is practiced. For this case, (7.95) goes to the following form:

$$\tau \approx \frac{\tau_{n0}}{2} \left(\frac{N_V}{N_C} \right)^{1/2} e^{[(E_C + E_V) / 2 - E_t] / kT} \quad (7.96)$$

The corresponding sphere in the $\ln\tau(1/T)$ curve is characterized by a decrease in the lifespan of unbalanced electrons and swells (Fig. 7.9-v, Sphere 3). With a further increase in temperature, when an inequality occurs in the area of private conductivity, $[(E_C + E_V) / 2 - E_t] \ll kT$ (7.95) can be written approximately as follows:

$$\tau = \frac{\tau_{p0}}{2} \left[1 + \left(\frac{N_C}{N_V} \right)^{1/2} \right] + \frac{\tau_{n0}}{2} \left[1 + \left(\frac{N_V}{N_C} \right)^{1/2} \right] \quad (7.97)$$

In this case, unbalanced electrons and hives tend to have a constant value. If $N_C = N_V$ that if accepted, This statement simplifies the following:

$$\tau = \tau_{p0} + \tau_{n0} = (\gamma_n + \gamma_p) / \gamma_n \gamma_p N_t \quad (7.98)$$

Thus, when the concentration of free charge carriers in the private conductivity region is very high with an increase in temperature, the residence time of non-equilibrium electrons and hives is determined only by the parameters of the recombination traps and is not dependent on temperature (Fig. 7.9- v, Fig. 4).

Knowing the concentration of electrons in saturation, it is possible to estimate E_t by the transition temperature from 1 sphere T_o , 2 spheres (Fig. 7.9- v). Based on the condition $T=T_o$, and $F=E_t$, we find:

$$n_0 = N_C e^{-(E_C - E_t)/kT_{0t}} \quad (7.99)$$

here

$$E_C - E_t = kT \ln(N_C / n_0) \quad (7.100)$$

2 Extrapolating the area of a straight line to the intersection with the coordinate axis, according to (7.93) $\ln(\tau_{p0} N_C / n_0)$ and find out t_0 from 1 sphere, $\ln(N_C / n_0)$ and put it in the formula (7.100). Also from the experimental correlation shown in Fig. 7.9 – $v \ln \tau(1/T)$ The depth of the surface of the recombination meat trap can be calculated.

7.9. Capture Centers and Recombination Traps

In the analysis of the recombination processes of charge carriers through traps, the trap electron and hives capture and the return of trapped electrons and hives from the traps to the corresponding zones under the influence of heat were taken into account. In this case, the phenomenon of thermal excitation of trapped charge carriers from the trap surface leads to an increase in residence time and, in turn, a decrease in the rate of recombination. Obviously, the intensity ratio of these processes depends on the state of the energy level of the traps.

Fig. 7.10 shows the energy scheme of a semiconductor containing local centers characterized by different levels in the exclusion zone. If there is a state other than thermodynamic equilibrium, then the concentration of free charge carriers is determined by the state of the quasi-Fermi surface, which is F_n for electrons and F_p for fumes. Therefore, the closer the energy level of the traps is near the bottom of the conduction zone, the greater the probability of thermal ionization of the charge carriers. Such centers are in a state of continuous electron exchange with the conduction zone, and do not contribute significantly to recombination processes. This can be called electron trapping or sticking traps. Near the upper edge of the tensile zone are the spurs of the pit snouting centers.

In traps that capture electrons and spikes, their recombination takes place, charge carriers are more likely to return to the corresponding zones due to heat, so recombination traps should be characterized by energy levels located deep enough in the lower part of the conduction zone.

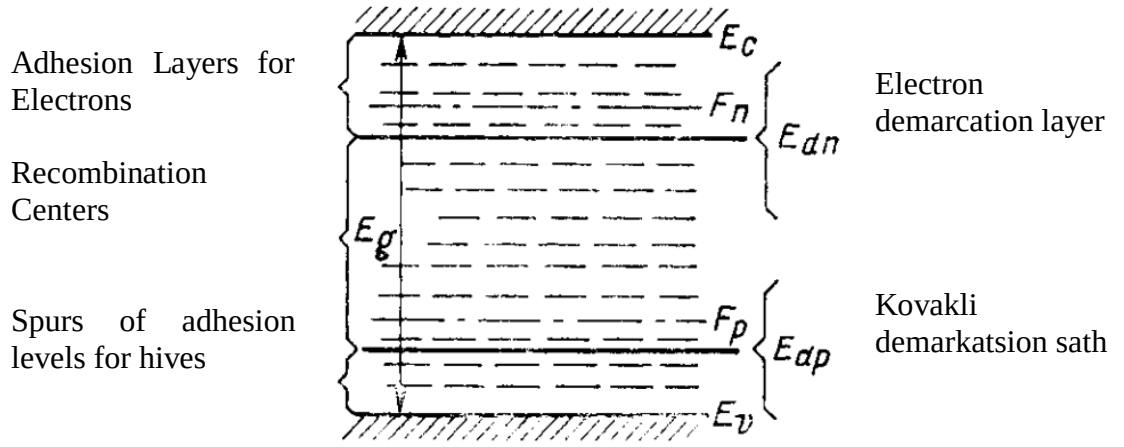


Fig. 7.10. Semiconductor containing capture centers and recombination traps

The relationship between the capture and recombination processes can be quantitatively described by introducing the coefficient k , which represents the ratio of the probability of capture of the opposite pointed charge carriers to the probability of heat return from the trap to the corresponding energy zone of the main charge carrier. With regard to electrons in the conduction zone, k_n is equal to the ratio of the probability of the spur being trapped by a negatively charged trap to the probability of an electron firing heat from the trap into the conduction zone. As can be seen from (7.70), the speed at which hives are caught by handles $\gamma_p N_t f_t p$ is determined by magnitude

The intensity of the heat excitation process of electrons from the handles according to (7.69) $\gamma_n N_t f_t n_1$ is equal to . here γ_n , γ_p the likelihood of electronic and hive catch, respectively f_t is an unbalanced distribution function for electrons located in traps; p is the concentration of unbalanced canines in the stress zone; n_1 is the equilibrium concentration of electrons in the conductivity zone when the Fermi level is equal to the trap level.

Taking this into account, we get the following expression for the k_n coefficient:

$$k_n = \frac{\gamma_p p}{\gamma_n n_1} = \frac{\gamma_p p}{\gamma_n N_C e^{-(E_C - E_t)/kT}} \quad (7.101)$$

Those centers where the probability of recombination is greater than the probability of separation due to the action of heat, i.e. $k_n > 1$, there will be recombination traps, and $k_n < 1$ traps are called catch centers. Energy level of the trap $k_n < 1$ Characterized by value, i.e., the recombination of charge carriers through traps and the probability of separation under the action of heat are the same, the level of circuit bounding (bounding) E_{dn} is called the Battle of the Bulge. Its status can be determined by the equality in the box:

$$\gamma_p N_V e^{(E_V - F_p)/kT} = \gamma_n N_C e^{-(E_C - E_{dn})/kT} \quad (7.102)$$

From here, we get the following:

$$E_{dn} = E_C + E_V - F_p - kT \ln \frac{\gamma_n N_C}{\gamma_p N_V} \quad (7.103)$$

From the attitude received, it follows that, Electronic arrivealanish of the level E_{dn} the state depends on the width of the prohibited zone, the position of the quasi-Fermi level of the hives, the effective mass of electrons and hives, the trap parameters (γ_n , γ_r) and depends on temperature. With an increase in the degree of excitation, which leads to an increase in the concentration of canines $F_p - E_V$ – Home The difference decreases, and the electron boundary surface approaches the lower part of the conductivity zone. In this case, the catch centers move into recombination traps. With an increase in temperature, the delimitation level moves into the middle of the forbidden zone, and the recombination traps move to the capture centers.

For recombination traps and capture centers whose surfaces are located near the upper edge of the tensile zone, the following coefficient may be introduced:

$$k_p = \frac{\gamma_n n}{\gamma_p p_1} \quad (7.104)$$

From the condition $k_n=1$, we find the state of the bounding level for the hives:

$$\gamma_n N_C e^{-(E_C - F_n)/kT} = \gamma_p N_V e^{(E_V - E_{dp})/kT} \quad (7.105)$$

or

$$E_{dp} = E_C + E_V - F_n - kT \ln \frac{\gamma_n N_C}{\gamma_p N_V} \quad (7.106)$$

Thus, if the energy level of the traps is higher than the electron boundary level, they will be the centers of electron reception. Traps with energy levels below the pit limit are the centers of hives catching. Among them is the energy level of the recombination traps (see Fig. 7.10). It should be noted that such a classification of the energy level of traps is rather conditional. In its implementation, in particular, the γ_n and γ_p of the capture coefficients changes weakly when the trap level changes. Capture centers can also contribute to recombination processes, which is especially noticeable in the absence of defects in the semiconductor that form energy levels in the middle part of the bandgap zone.

Control Questions

1. Explain the generation of electrons and cavities
2. How does the recombination of electrons and spurs occur?
3. Explain the recombination mechanisms
4. What do you mean by interband radiation recombination?
5. How does the recombination of charge carriers by traps occur?
6. Give Insights on Balanced and Unbalanced Chargers

7. Explain bipolar optical generation of charge carriers
8. How does the monopolar optical generation of charge carriers occur?
9. What does Maxwell relativity time mean?

Chapter VIII.

DIFFUSION AND DRIFT OF UNBALANCED CHARGE CARRIERS

8.1. Continuity equation

A semiconductor may have non-neuivial charge carriers due to external influences. In this case, the $n(x, y, z, t)$ $p(x, y, z, t)$ function of the spatial coordinates and the function of time is characterized by the continuity of the rate of their change.

Let us cite the continuity equation for a one-dimensional, homogeneous specimen, in which the continuity of the electrons change only in the direction of its x axis (Fig. 8.1).

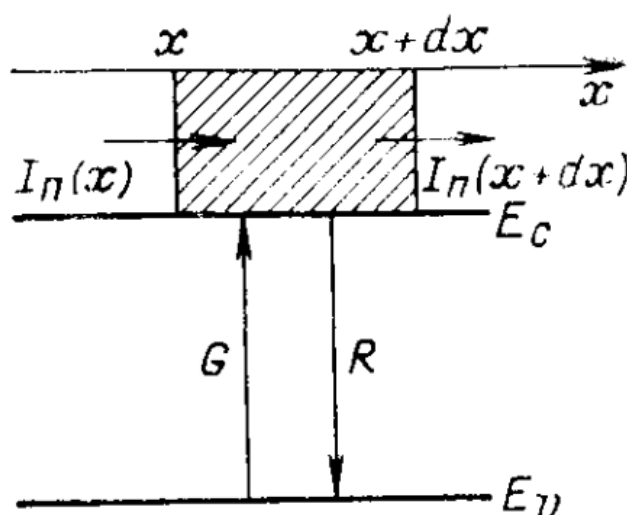


Fig. 8.1. Plot to explain the continuity equation

In a semiconductor $1\text{cm}^2 \cdot dx$, we distinguish a layer with a thickness. This layer dx is equal to thickness. Suppose that the number of electrons in time t is $n(x,t)$, then the number of electrons in this time period is $n(x,t)dx$. Change in the number of electrons in the time period $t+dt$ $n(x,t+dt)$ and dt in time time dx :

$$n(x, t + dt)dx - n(x, t)dx = \frac{\partial n}{\partial t} dx dt \quad (8.1)$$

This is due to the change in the number of electrons, recombination, as well as the dredging and diffusion of charge carriers. We see these processes:

1. Electron heinsion - let the semiconductor be extinguished as a result of the absorption of light from the beam and be equal to the electroctron-covak pair that is being decomposed at $1\text{cm}^3 \cdot 1 \text{ s } G$. He $dx \cdot dt$ was in the middle of a time in the Hall of Fame

$$G dx dt \quad (8.2)$$

And the electorates are the ones who make the difference.

2. dx Changes in the number of charge carriers in the hajm can be caused by recombination. The full speed of recombination $R_0 + R$ is here, R_0 – the rate of recombination in the heat mode, the rate of recombination in R the presence of external influences. In the case of linear recombination, the change in the number of electrons due to the recombination of R of dt in dt time, according to (7.18), if the lifetime of the non-equilibrium electrons does not depend on the electron concentration τ_n , the change in the number of electrons due to the recombination of dx in time dx is as follows:

$$Rdxdt = -\left(\frac{dn}{dt}\right)_r dxdt = \frac{n - n_0}{\tau_n} dxdt \quad (8.3)$$

3. Changes in the amount of charge carriers in dx are caused by diffusion and drainage processes. If $I_n(x,t) \cdot 1 \text{ sm}^2$ is a stream of electrons passing through 1 s , then let dt be the number of electrons entering the dx layer boundary from the dt layer boundary, $I_n(x,t)dx$, $x + dx$ the number of electrons exiting the layer boundary $I_n(x+dx,t)dx$. As a result, due to this flux difference, the change in the number of electrons at time dx is as follows:

$$I_n(x,t)dt - I_n(x+dx,t)dt = -\frac{\partial I_n}{\partial x} dxdt \quad (8.4)$$

A complete change in the number of electrons during dt time in dx :

$$\frac{\partial n}{\partial t} dxdt = Gdxdt - \frac{\partial I_n}{\partial x} dxdt - \frac{n - n_0}{\tau_n} dxdt \quad (8.5)$$

From here,

$$\frac{\partial n}{\partial t} = G - \frac{\partial I_n}{\partial x} - \frac{n - n_0}{\tau_n} \quad (8.6)$$

(8.6) The equation is called the continuity equation for electrons. Similarly, the continuity equation for puppies is written as follows:

$$\frac{\partial p}{\partial t} = G - \frac{\partial I_p}{\partial x} - \frac{p - p_0}{\tau_p} \quad (8.7)$$

Taking into account the charge carriers, we record the currents of electrons I_n and particles I_p , through the densities of the charges, as follows:

$$j_n = -eI_n, \quad j_p = eI_p \quad (8.8)$$

If the continuity of charge carriers (x,y,z) is a function of the coordinates, the continuity tension is written through the current densities as follows:

$$\frac{\partial n}{\partial t} = G + \frac{1}{e} \text{div} j_n - \frac{n - n_0}{\tau_n} \quad (8.9)$$

$$\frac{\partial p}{\partial t} = G - \frac{1}{e} \text{div} j_p - \frac{p - p_0}{\tau_p} \quad (8.10)$$

Here,

$$\text{div} \mathbf{j} = \frac{\partial j_x}{\partial x} + \frac{\partial j_y}{\partial y} + \frac{\partial j_z}{\partial z}.$$

Let there be a source of non-ambiguous charge carriers at point $x=0$. Consider the stationary status. In a stationary state, the continuity of the electrons and covalents does not change over time, i.e. $\frac{\partial n}{\partial t} = \frac{\partial p}{\partial t} = 0$., in a one-dimensional holo, the continuity equation is as follows:

$$-\frac{1}{e} \frac{\partial j_n}{\partial x} = G - \frac{n - n_0}{\tau_n} \quad (8.11)$$

$$\frac{1}{e} \frac{\partial j_p}{\partial x} = G - \frac{p - p_0}{\tau_p} \quad (8.12)$$

Equations (8.11) and (8.12) represent the conditions under which the number of charge carriers can be conserved. Consequently, at steady state, without taking into account recombined charge carriers in volume, the flow of charge carriers flowing out of the volume becomes equal to the number of charge carriers formed by external excitation.

8.2. Diffusion and diffusion

Now consider a non-homogeneous semiconductor in which the concentration of electrons $n(\mathbf{r})$ and $p(\mathbf{r})$ varies from that point to that point. Therefore, a diffusion flux of electrons and spikes must occur in the non-homogeneous semiconductor, which can be determined by the scattering of charge carriers from the regions with a high concentration to the places with a low concentration.

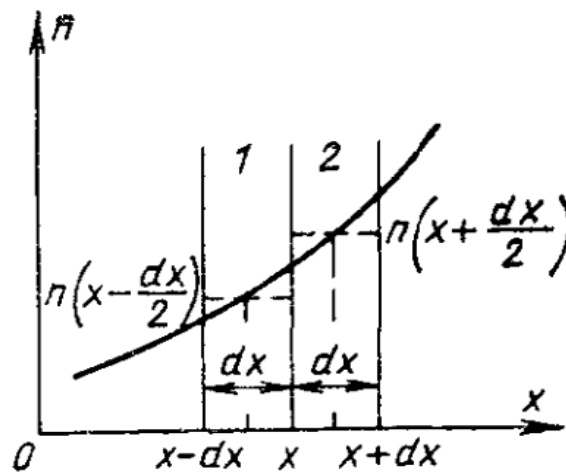


Fig. 8.2. Graph for estimating the concentration gradient of charge carriers

Suppose that in a semiconductor, the density of charge carriers in the direction of the x -axis is increased (Fig. 8.2). We conduct a perpendicular investigation of the x axis of the point x , and from this axis we see the behavior of the charge carriers in layers 1 and 2 of thickness dx , which are aligned on the left and right axis. As a result of the random characterization, one part of the charge carriers is divided from layer 1 to layer 2, and a part of the fifth is divided from layer 2 to layer 1. The number of charge carriers in layer 2 is greater than in layer 1, where the current of charge carriers is greater than the current in the right direction.

If $n(x - \frac{dx}{2})$ there is an average concentration of electrons in -1 layer, electrons in $n(x + \frac{dx}{2})$ -2 layers, then the difference in electrification in these layers is as follows:

$$n(x - \frac{dx}{2}) - n(x + \frac{dx}{2}) = -\frac{dn}{dx} dx \quad (8.13)$$

According to (8.13), the difference I_n , concentration of electrons is proportional to the gradient of their concentration, so that the flux of electrons resulting from their diffusion along the direction of their x -axis is in proportional to the gradient of concentration of electrons in that direction. The flux of electrons in the direction of the x -axis resulting from diffusion can be written as follows:

$$I_n = -D_n \frac{dn}{dx} \quad (8.14)$$

Here, D_n the diffusion coefficient of -electrons.

Similar Puppets Text:

$$I_p = -D_p \frac{dp}{dx} \quad (8.15)$$

is equal to . D_p -Diffusion of Molds.

From Fig. (8.14) and (8.15) it is clear that the axes of the electrons and the coves are tilted towards the tómon, where the coffin of the charge carriers is small.

Diffusion shocks caused by diffusion of charge carriers include:

$$j_{nDif} = eD_n \frac{dn}{dx} \quad (8.16)$$

$$j_{pDif} = eD_p \frac{dp}{dx} \quad (8.17)$$

If n the functions of p the and (x,y,z) coordinates are calculated, the diffusion currents are in vector form, for electrons

$$j_{nDif} = eD_n \text{grad}n(r) \quad (8.18)$$

For hives

$$j_{pDif} = eD_p \text{grad}p(r) \quad (8.19)$$

So so.

The diffusion current, which is caused by the presence of a concentration gradient of the charge carriers, leads to spatial splitting of the charges and to the emergence of a static electric field and produces flux currents of electrons and cavities. In thermodynamic equilibrium, the drift current at each point in the semiconductor equalizes the diffusion current, so that the total current becomes zero.

Suppose that a homogeneous semiconductor field is located in an external constant electric field whose field strength is E . Under the influence of this field, electrons and spikes will have a directed motion, resulting in an electron and a transmissible current. If the external electric field is weak and does not change the character of the motion of the charge carriers, then the drift organizer of the current density is written according to Ohm's law as follows:

$$j_{ndr} = en\mu_n\mathcal{E}; \quad j_{pdr} = ep\mu_p\mathcal{E}. \quad (8.20)$$

A full tök is a collection of diffusion and dréyf tök. We write full expressions for electrifications and puppies:

$$j_n = j_{ndr} + j_{nDif} = en\mu_n\mathcal{E} + eD_n \frac{dn}{dx} \quad (8.21)$$

$$j_p = j_{pdr} + j_{pDif} = ep\mu_p\mathcal{E} - eD_p \frac{dp}{dx} \quad (8.22)$$

Thus, the total density of a non-homogeneous semiconductor at any point and at any moment of time is as follows:

$$j = j_n + j_p = e(n\mu_n + p\mu_p)\mathcal{E} + e(D_n \frac{dn}{dx} - D_p \frac{dp}{dx}) \quad (8.23)$$

It is important to note that the diffusion current is only relevant in semiconductors. This is because while the concentration of all aggregate charges in semiconductors is unchanged, the concentrations of electrons and spikes can vary over a wide range. In metals, the concentration of electrons is almost unchanged.

8.3. Einstein's Relationship

In a non-homosexual semiconductor, at thermodynamic equilibrium, the current is zero, that is, $J = J_n + J_p = 0$. In this case, the conductivity current is able to equilibrium the diffusion current and, based on (8.21), for electrons we can write:

$$n\mu_n E_{doim} = -D_n \frac{dn}{dx} \quad (8.24)$$

$E_{permenant}$ Since a semiconductor has a static electric field, the electrons located in that field $U = -e\varphi$ will have potential energy. Therefore, when the semiconductor is not deflected, the concentration of the electron in the conduction band satisfies the Boltzman attitude.

$$n = N_C e^{-(E_C + U - F)/kT} = n_0 e^{e\varphi/kT} \quad (8.25)$$

here

$$n_0 = N_C e^{-(E_C - F)/kT}$$

- Concentration of electrons in equilibrium; φ is the electrostatic potential.

$$E_{doim} = -\frac{d\varphi}{dx}$$

Considering that n and $\frac{dn}{dx}$ If we put the values of (8.24) into the equation, we get:

$$-\mu_n n_0 e^{e\varphi/kT} \frac{d\varphi}{dx} = -D_n \frac{en_0}{kT} e^{e\varphi/kT} \frac{d\varphi}{dx} \quad (8.26)$$

wherein for electrons:

$$\frac{\mu_n}{D_n} = \frac{e}{kT} \quad (8.27)$$

The same is true for holes:

$$\frac{\mu_p}{D_p} = \frac{e}{kT} \quad (8.28)$$

The equation that connects the diffusion coefficient of charge carriers subject to Maxwell's statistics to their drift mobility under thermodynamic equilibrium conditions is called the Einstein relation (Einstein relativity).

Experience has shown that Einstein's attitude also applies to unbalanced charge carriers. This is natural, because unbalanced charge carriers come into heat equilibrium with the grid by exchanging energy with the grid in a much shorter period of time than their residence time, and in the absence of equilibrium, their energy distribution is no different from the distribution of charge carriers in equilibrium.

8.4. Diffusion and drift of unbalanced charge carriers in monopolar conductivity

In the case of monopolar conductivity, let us consider the diffusion and drift of unbalanced charge carriers, in which free charge carriers arise only as a result of their excitation from their intake surface. Suppose a part of a sufficiently long monochemical semiconductor, for example, of the n-type, is illuminated by weak absorbing light. A homosexual generation of electrons in the illuminated part of a semiconductor at $x < 0$ (Fig. 8.3 - a) occurs as a result of their seeding from the donor state to the conduction zone under the influence of light.

In this case, the concentration of non-random electrons is determined by the concentrations of n_0 and excess electrons n in equilibrium excited from the n donor

input level Δn . If there are no capture centers, then the number of excess electrons Δn is equal to the number of donor input positive ions ΔN_d^+ , i.e.

$$\Delta n = \Delta N_d^+ \quad (8.29)$$

This equation represents the condition of electric neutrality in the case of monopolar conductivity.

Since the concentration of electrons in the illuminated part of the sample is higher than in the non-luminous part, the irbalanced electrons in the illuminated part of the sample are diffused into the non-luminous part (Fig. 8.3– b). As a result, electric neutrality is violated in a certain region of the semiconductor, resulting in a volumetric charge as well as an electric field.

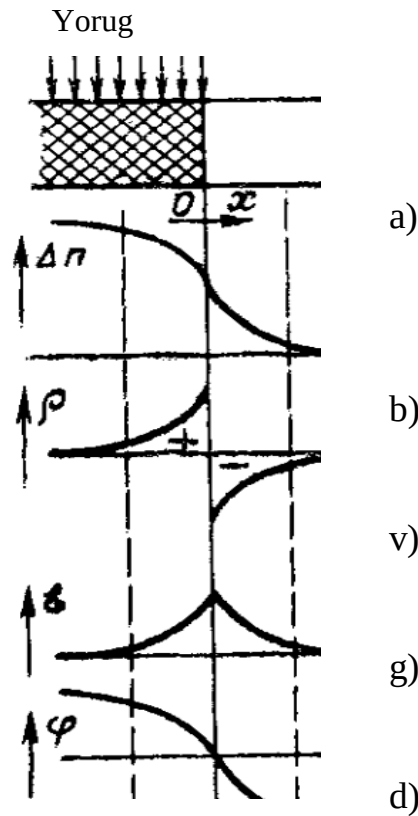


Fig. 8.3. The distribution of the donor semiconductor is (a) concentration in partial illumination (b), volumetric charge (v), electric field (g), potential (d)

As a result of diffusion, the volumetric charge ρ is negative in the unlit part of the sample from which the excess electrons come, while in the illuminated area it is positive at the expense of the ions of the donor inputs where they left (Fig. 8.3-v). These charges generate a static electric field due to the directed E_{bar} voltage, which is directed in such a way that it prevents the diffusion of unbalanced electrons (Fig. 8.3-g).

Thus, the occurrence of diffusion current $J_{dif} = eD_n dn/dx$ results in the generation of a static electric field having a voltage E and resulting in a directed conductivity current against the diffusion current $J_{o't} = en\mu_n E_{bar}$. In steady state, the density of full current is zero:

$$J = J_{o't} + J_{dif} = en\mu_n E_{bar} + eD_n dn/dx = 0 \quad (8.30)$$

From (8.30) it is possible to determine the strength of the static electric field. For the low level of optical excitation, we calculate E_{bar} when the concentration of excess electrons is small relative to its concentration in equilibrium, i.e., when $\Delta n \ll n_0$. Using Einstein's attitude, we get:

$$E_{bar} = -\frac{D_n}{n\mu_n} \frac{d\Delta n}{dx} = -\frac{kT}{e(n_0 + \Delta n)} \frac{d\Delta n}{dx} \approx -\frac{kT}{en_0} \frac{d\Delta n}{dx} \quad (8.31)$$

Differentiating (8.31) we get:

$$\frac{dE_{bar}}{dx} = \frac{kT}{en_0} \frac{d^2 \Delta n}{dx^2} \quad (8.32)$$

$\frac{dE_{bar}}{dx}$ The value can be found using the Poisson equation:

$$\frac{dE_{bar}}{dx} = \frac{1}{\epsilon_r \epsilon_0} \rho = -\frac{e}{\epsilon_r \epsilon_0} \Delta n, \quad (8.33)$$

Here $\rho = -e\Delta n$ is the volumetric charge density in the unlit part of the sample. From equations (8.32) and (8.33) follows the following conclusion:

$$\frac{d^2 \Delta n}{dx^2} - \frac{e^2 n_0}{\epsilon_r \epsilon_0 kT} \Delta n = 0 \quad (8.34)$$

If the $l_e = \sqrt{\epsilon_r \epsilon_0 kT / e^2 n_0}$ value is then (8.34) is written as follows:

$$\frac{d^2 \Delta n}{dx^2} - \frac{\Delta n}{l_e^2} = 0 \quad (8.35)$$

Here's a general solution to the equation:

$$\Delta n = C_1 e^{\alpha_1 x} + C_2 e^{\alpha_2 x}, \quad (8.36)$$

where C_1 and C_2 are constants determined from boundary conditions; α_1, α_2 are the roots of the characteristic equation, which are equal to:

$$\alpha_1 = \frac{1}{l_e}; \quad \alpha_2 = -\frac{1}{l_e}; \quad (8.37)$$

The concentration of excess electrons in the unlit region of the semiconductor decreases as it moves away from the illuminated part of the sample, only the solution of the case with a degree negative indicator makes sense, so

$$\Delta n = \Delta n(0) e^{-x/l_e}. \quad (8.38)$$

Thus, in the case of monopolar conductivity, the concentration of excess charge carriers in the unlit part of the sample decreases exponentially with the

decreasing constant l_e with distance from the illuminated area, and this is called the shielding radius (length) or Debye radius. The shielding length, according to (8.35), depends on the concentration of the main charge carrier, so its value can vary over wide ranges depending on the semiconductor's specific conductivity. For example, for non-decayed semiconductors, such as germanium and silicon, the shielding radius can be 10^{-4} - 10^{-6} cm, depending on the degree of alloying.

From the formula (7.29), it follows that the volumetric charge introduced into the semiconductor will be present for an average time t as a result of the conduction current after the excitatory light is turned off.

$$\rho = \rho_0 e^{-t/\tau_\mu}. \quad (8.40)$$

If the density of a volume charge divided by unit charge, we get the change in the concentration of excess charge carriers over time:

$$\Delta n = \Delta n(0) e^{-t/\tau_\mu}. \quad (8.41)$$

From a comparison of expressions in (8.39) and (8.41), it follows that in the monopolar state, the propagation of charge carriers over the shielding l_e length occurs during the Maxwell relaxation time τ_μ , which is the effective time for establishing a diffusion-drift equilibrium.

8.5. Diffusion and drift of non-basic excess charge carriers in input semiconductors

Let us consider the diffusion and drift of non-primary overcharge carriers in an input semiconductor when there is a homogeneous external electric field whose voltage is E .

Suppose there is an n-type semiconductor ($p \ll n_0$) the transverse dimensions of which are significantly smaller than its length. A very narrow region of a semiconductor ($-l \leq x \leq 0$) is illuminated by light, so that a uniform generation of electrons and spikes occurs over the entire volume of the sample's illuminated region (Fig. 8.4). Let the concentration of the principal charge carriers be greater relative to the concentration of non-primary charge carriers p , where the residence time Δp is n_0 .

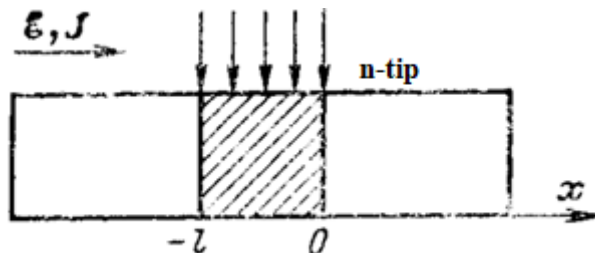


Fig. 8.4. One-dimensional model of semiconductor sample

The concentration of non-equilibrium holes of an electron semiconductor with $\Delta n = \Delta p$ an equal amount of excess electrons and holes in the illuminated field $\Delta n = p$ is significantly greater than the concentration of p_0 Δ in the unlit part of the sample $p = p_0 + \Delta p$. Due to the presence of a concentration gradient of the holes, their diffusion occurs. In the unlit part of the sample, where the holes are diffused, a volumetric positive charge is generated and the electrical neutrality of the sample is violated. This volumetric charge generates a static electric field that causes the charge carriers to redistribute, so that during Maxwell relaxation, electrons are attracted from the illuminated part of the sample to that region, and the charge of the excess holes is compensated. If the concentration of excess electrons is equal to the concentration of excess holes at any point in the sample, a full compensation of the volumetric charge occurs. Thus, excess non-basic charge carriers—holes, which have an equal number of primary charge carriers with them when deeply diffused into an electron semiconductor—carry electrons, so that no volumetric charge is generated, and their diffusion occurs as a quasi-neutral cloud diffusion of non-equilibrium carriers with charge $\Delta n = \Delta p$. As it penetrates deeper into the semiconductor, the excess holes and electrons recombine, and their concentration decreases with distance.

If such a semiconductor is exposed to a homogeneous electric field E , then in the accelerating field the holes will attract electrons by an equal amount with them, and if the field slows down the movement of the holes, then a corresponding redistribution of the main charge carriers will occur.

Consequently, in the external electric field, the quasi-neutral cloud of unbalanced charge carriers in the n-type material moves at the drift velocity of non-primary carriers in the direction of motion of the excess holes.

The concentration of excess holes Δ can be determined by solving the continuity equation p :

$$\frac{\partial p}{\partial t} = G - \frac{1}{e} \frac{\partial J_p}{\partial x} - \frac{p - p_0}{\tau_p} \quad (8.42)$$

Let us solve this equation for the stationary state in the unlit part ($G=0$) $\frac{\partial p}{\partial t} = 0$ of the semiconductor. To do this, by differentiating the density of the hole current (8.22) and putting it in (8.42), we get the following:

$$D_p \frac{d^2 \Delta p}{dx^2} - \mu_p E \frac{d \Delta p}{dx} - \frac{\Delta p}{\tau_p} = 0 \quad (8.43)$$

This equation is valid for an external electric field in such a way that D_p , μ_p , τ_p should not be dependent on the field.

Breaking down the two parts of (8.43) by D_p and following

$$L_p = \sqrt{D_p \tau_p} \quad (8.44)$$

$$L_E = \tau_p \mu_p E, \quad (8.45)$$

By introducing the definition (8.43), we make it look like this:

$$\frac{d^2 \Delta p}{dx^2} - \frac{L_E}{L_p^2} \frac{d \Delta p}{dx} - \frac{\Delta p}{L_p^2} = 0 \quad (8.46)$$

The overview of this equation would be as follows:

$$\Delta p = C_1 e^{\alpha_1 x} + C_2 e^{\alpha_2 x}, \quad (8.47)$$

where C_1 and C_2 are constants determined from boundary conditions;

$$\Delta p|_{x=\Delta p}|_{x=-l} = \Delta p(0), \quad x \rightarrow \pm\infty \quad \Delta p \rightarrow 0; \quad \alpha_1, \alpha_2 -$$

The roots of the characteristic equation, which are equal to:

$$\alpha^2 - \frac{L_E}{L_p^2} \alpha - \frac{1}{L_p^2} = 0 \quad (8.48)$$

$$\alpha_{1,2} = \frac{L_E \pm \sqrt{L_E^2 + 4L_p^2}}{2L_p^2} \quad (8.49)$$

Considering the decrease in the concentration of non-primary non-equilibrium charge carriers as the sample moves away from the illuminated area, we finally get the following: $x \geq 0$

$$\Delta p = C_1 e^{-x/L_1} = \Delta p(0) e^{-x/L_1} \quad (8.50)$$

here

$$L_1 = \frac{2L_p^2}{\sqrt{L_E^2 + 4L_p^2} - L_E} \quad (8.51)$$

and

$$\Delta p = C_2 e^{x/L_2} = \Delta p(0) e^{x/L_2} \quad (8.52)$$

here

$$L_2 = \frac{2L_p^2}{\sqrt{L_E^2 + 4L_p^2} + L_E} \quad (8.53)$$

Thus, the concentration of excess non-primary charge carriers on either side of the sample's illuminated region decreases exponentially on the basis of the L_1 and L_2 reduction constants, which are referred to as the gravity length.

Let us analyze the resulting expressions. Let there be no external electric field ($E=0$). In the absence of an external electric field, only the diffusion of excess canopies occurs. $L_E = \tau_p \mu_p = 0$ Since the excess concentration of canines of non-main charge carriers as a result of recombination varies with distance, according to the following law:

$$\Delta p = \Delta p(0)e^{-x/L_p}. \quad (8.54)$$

It follows from this expression that the exponential decrease in the concentration of excess non-primary charge carriers caused by recombination is symmetrical on both sides of the illuminated portion of the sample (Fig. 8.5-a), and L_p is determined by the value of the non-primary charge carriers here, called the diffusion length of the spurs. L_p is the mean distance of displacement of unbalanced hives during diffusion. Therefore, in the absence of an external electric field, the concentration of excess canines as a result of recombination changes exponentially with the law of distance, according to (8.54). This distribution also applies to non-equilibrium electrons, since it is hypothesized that electron neutrality will occur when solving the problem, according to which $\Delta n = \Delta p$.

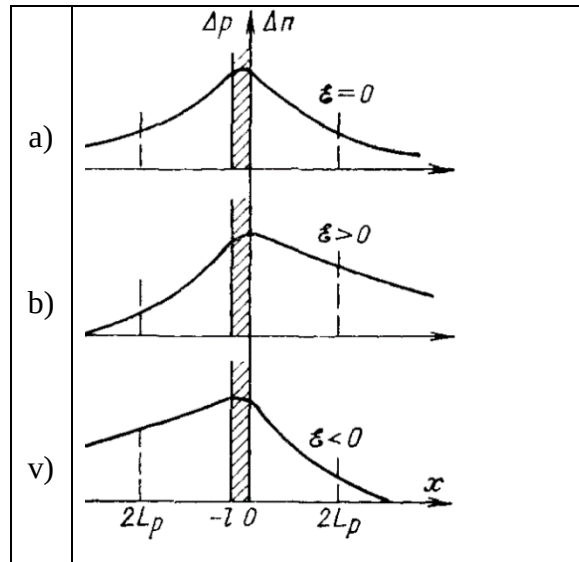


Fig. 8.5. Distribution of Concentration of Excess Charge Carriers Across the Sample

Since only the dispersion of excessive charge carriers occurs, the current is a diffusion current and becomes as follows:

$$J_{p\,dif} = -eD_p \frac{dp}{dx} = \frac{eD_p}{L_p} \Delta p = \frac{eL_p}{\tau_p} \Delta p = J_{p\,dif}(0)e^{-x/L_p}, \quad (8.55)$$

That is, the diffusion current of surplus charge carriers changes with the coordinate in accordance with the changed law in which the concentration of surplus fluids is changed.

$$v_D = \frac{L_p}{\tau_p} = \frac{D_p}{L_p}, \quad (8.56)$$

The magnitude is called the diffusion rate. In numerical terms, it is equal to the rate at which unbalanced hives travel a path, which is equal to the length of diffusion during their lifetime.

Now $E \neq 0$ let's look at the situation. If $L_E \ll 2L_p$, or $eL_p E \ll 2kT$ (6.183), then $L_1 = L_2 = L_p$. Consequently, when $E < E_c$ in a weak electric field $E \ll E_c$, where $E_c = 2kT / eL_p$ is the critical field, the distribution of the excess concentration of non-basic charge carriers is determined by diffusion and is characterized by Equation (8.54).

In the presence of a strong external electric field, when $|L_E| > 2L_p$, the reduction constants L_1 and L_2 diffusion length vary from L_p . Depending on the direction of the electric field ($E > 0$ and $E < 0$), they are greater or smaller than L_p . For example, in $E > 0$ for the unlit part of the sphere, $x < -l$ and $x > 0$,

$$L_2 = \frac{2L_p^2}{\sqrt{L_E^2 + 4L_p^2} + L_E} \langle L_p \rangle = \frac{2L_p^2}{\sqrt{L_E^2 + 4L_p^2} - L_E} = L_1. \quad (8.57)$$

It follows that the external electric field breaks the symmetry in the distribution of the concentration of excess charge carriers.

We now $L_E^2 \gg 4L_p^2$ consider the case of such large areas in which inequality $E \gg E_c$ is satisfied. This condition can be written in accordance with (8.45) and (8.56) as follows:

$$\frac{L_E^2}{4L_p^2} = \frac{(\tau_p \mu_p E)^2}{4L_p^2} = \frac{\tau_p^2 v^2}{4L_0^2} = \frac{v^2}{4v_D^2} \gg 1, \quad (8.58)$$

Here $v = \mu_p E$ -drift speed.

The distance traveled by non-equilibrium charge carriers at the rate of drift during their lifetime is called the length of the drift of magnitude L_E , which is numerically equal to it. Therefore, if the length of drift is much greater than the length of the diffusion, the field voltage E is also large.

We define the lengths L_1 va L_2 in the case that $E > 0$. From the expression for the sphere of $X > 0$ (8.51) it is possible to write the following:

$$L_1 = \frac{2L_p^2}{L_E \sqrt{1 + 4L_p^2 / L_E^2} - 1} \approx \frac{2L_p^2}{L_E} \frac{1}{2L_p^2 / L_E^2} = L_E, \quad (8.59)$$

That is, the length of gravity is equal to the length of the drift

$$L_1 = L_E = \tau_p \mu_p E = \tau_p v \quad (8.60)$$

and the distribution of the concentration of excess carriers is determined by the following expression:

$$\Delta p_1 = \Delta p(0) e^{-x/L_1} = \Delta p(0) e^{-x/\tau_p \mu_p E}, \quad (8.61)$$

That is, just as in the diffusion-only case, the concentration of excess hives decreases exponentially with increase x , but now the reduction will be constant $L_1 = L_E$. Since the reduction is constant $L_1 > L_p$, it L_1 is also called "length of diffusion across the field".

From comparisons of expressions (8.54) and (8.61), it follows that $L_E^2 > 4L_p^2$ the excess concentration of charge carriers in the $x > 0$ region of the semiconductor is greater than in the absence of an external electric field (Fig. 8.5-b).

Consequently, when the current emanating from a strong electric field passes, i.e., when the rate of drift is much greater than the diffusion rate ($v \gg v_D$), when $E > 0$, the excess spurs in the electronic semiconductor are attracted by the field towards the $x > 0$ region of the semiconductor, and the semiconductor is enriched with non-basic charge carriers in the absence of an external electric field only in the presence of diffusion. This phenomenon is called non-primary charge carrier injection.

For a P-type semiconductor, electron injection is observed at $E < 0$. Now we consider the distribution of excess hives for the $x < -l$ sphere, $E > 0$ case. Now the length of compression

$$L_2 = \frac{2L_p^2}{L_E \sqrt{1 + 4L_p^2/L_E^2 + 1}} \approx \frac{2L_p^2}{L_E} \frac{1}{2 + 2L_p^2/L_E^2} \approx \frac{L_p^2}{L_E} \quad (8.62)$$

and based on (8.52), we get:

$$\Delta p_2 = \Delta p(0) e^{-x/L_2} = \Delta p(0) e^{-L_E/L_p^2}. \quad (8.63)$$

It can be seen that as the field strength increases, the L_E increases, while L_2 decreases. If L_2 is $< L_p$, and for all points of the semiconductor in the region $x < -l$, e.g. $2L_p \Delta p_2 < \Delta p$ at a distance of L_2 , is called the "diffusion length against the field". This means that as the electric field strength increases at $E > 0$, non-primary charge carriers in the $x < -l$ region of the electron semiconductor volume become poorer (left side of Fig. 8.5-b). This phenomenon is termed non-primary charge carrier exclusivity.

In a P-type semiconductor, the electron expression occurs at $E < 0$. When the direction of the external electric field $E < 0$ changes, there is a decrease in the concentration of excess charge carriers in the region $x > 0$ of the semiconductor, and their increase in the region $x < -l$ (Fig. 8.5-v). These phenomena are respectively

termed extraction and accumulation of unequilibrium charge carriers. In a p-type semiconductor, they are observed at $E > 0$.

8.6. Diffusion and drift of unbalanced charge carriers in a semiconductor near private conductivity

Let us consider the diffusion and drift of non-equilibrium charge carriers in a semiconductor where the number of electrons and spikes in equilibrium is equal, that is, in a semiconductor whose conductivity is close to the individual. We assume that unbalanced charge carriers are triggered by the transition from zone to zone. If the electrons and hives are not captured by the traps, then $\Delta n = \Delta p$.

Suppose that in such a homogeneous semiconductor there is a non-homosexual and distribution of unbalanced charge carriers resulting from the local illumination of the sample along the x-axis $n(x) = n_0 + \Delta n(x)$ $p(x) = p_0 + \Delta p(x)$. We consider the level of excitation to be low, in which $\Delta n = \Delta p \ll (n_0 + p_0)$.

Since there is an external electric field and a concentration gradient of the charge carriers, diffusion and drift currents of electrons and cavities occur, which are characterized by (8.21) and (8.22), and the current density is equal to:

$$J = J_n + J_p = e(n\mu_n + p\mu_p)E + e(D_n dn/dx - D_p dp/dx). \quad (8.64)$$

In an insulated semiconductor with a stationary state, the total current is zero ($J = 0$). In this case, at each point of the sample, the diffusion currents are balanced by the drift currents arising from a static electric field of E_{bar} voltage, which we find by assuming $J = 0$ (8.64):

$$E_{bar} = \frac{D_p \frac{dp}{dx} - D_n \frac{dn}{dx}}{n\mu_n + p\mu_p} = \frac{D_p \frac{dp}{dx} - D_n \frac{dn}{dx}}{(n_0 + \Delta n)\mu_n + (p_0 + \Delta p)\mu_p}. \quad (8.65)$$

It follows from (8.65) that the smaller the static electric field E_{st} generated by the diffusion of charge carriers in the absence of electric current ($J=0$), the greater the concentrations of electrons n_0 and the p_0 equilibrium concentrations of the spurs. If the specific electrical conductivity of a semiconductor is high enough, and the concentration of excess charge carriers Δn and Δp is significantly less than the concentration of charge carriers in equilibrium Δn and p , i.e. $\Delta n \ll n_0$ and $\Delta p \ll p_0$. If so, then it can be assumed that there is no static electric field at the first approximation, and the electroneutrality condition is provided by the fact that $n \Delta = \Delta p$ at each point of the semiconductor. Consequently, the diffusionable charge carriers carry opposite pointed charge carriers equal during the diffusion. In this case, the process of

diffusion of excess electrons and cavities occurs as freely as the diffusion of excessively neutral electron pit pairs, characterized by the same residence time τ .

Suppose that the sample generated by the unbalanced electron-cavity pairs is significantly larger than the internal static field, i.e., $E \gg E_{st}$ is located in the outer homogeneous electric field E . For simplicity's sake, we limit ourselves to a one-dimensional state and assume that the concentration gradient and the external electric field are directed along the x-axis. In this case, the continuity equations and the current density equation should be written for electrons and swells as follows:

$$\left. \begin{aligned} \frac{\partial \Delta n}{\partial t} &= \frac{1}{e} \frac{\partial J_n}{\partial x} - \frac{\Delta n}{\tau}; \\ \frac{\partial \Delta p}{\partial t} &= -\frac{1}{e} \frac{\partial J_p}{\partial x} - \frac{\Delta p}{\tau}; \end{aligned} \right\} \quad (8.66)$$

$$\left. \begin{aligned} J_n &= \sigma_n E + e D_n \frac{\partial \Delta n}{\partial x}; \\ J_p &= \sigma_p E - e D_p \frac{\partial \Delta p}{\partial x}, \end{aligned} \right\} \quad (8.67)$$

here $\sigma_n = en\mu_n$ and $\sigma_p = ep\mu_p$ the electronic and pit constituents of specific conductivity.

Under the influence of an external electric field, a constant velocity drift of electron-hive pairs is observed. wherein the sample is characterized by the combined diffusion and drift of electrons and hives, provided that the sample is electron-neutral, with an effective drift mobility μ_E and an effective diffusion coefficient D , which is the same for electrons and hives. μ_E magnitudes E and D are also known as ambiquity drift mobility and ambiquity diffusion coefficients.

μ_E determine magnitudes E and D , we write the continuity equations (8.66), we put the values of J_n and J_p from (8.67) on them:

$$\left. \begin{aligned} \frac{\partial \Delta n}{\partial t} &= D_n \frac{\partial^2 \Delta n}{\partial x^2} + \mu_n E \frac{\partial \Delta n}{\partial x} - \frac{\Delta n}{\tau}; \\ \frac{\partial \Delta p}{\partial t} &= D_p \frac{\partial^2 \Delta p}{\partial x^2} - \mu_p E \frac{\partial \Delta p}{\partial x} - \frac{\Delta p}{\tau}; \end{aligned} \right\} \quad (8.68)$$

We multiply (8.68) by σ_p and σ_n , respectively, and add both equations. As a result, taking into account $\Delta p = \Delta n$ and using Einstein's attitude, we get:

$$\frac{\partial \Delta n}{\partial t} = \frac{D_n \sigma_p + D_p \sigma_n}{\sigma_n + \sigma_p} \frac{\partial^2 \Delta n}{\partial x^2} + \frac{\mu_n \sigma_p - \mu_p \sigma_n}{\sigma_n + \sigma_p} E \frac{\partial \Delta n}{\partial x} - \frac{\Delta n}{\tau} \quad (8.69)$$

For the stationary state, i.e., $\frac{\partial \Delta n}{\partial t} = 0$ when (8.69) can be written as follows:

$$\frac{D_n \sigma_p + D_p \sigma_n}{\sigma_n + \sigma_p} \frac{\partial^2 \Delta n}{\partial x^2} + \frac{\mu_n \sigma_p - \mu_p \sigma_n}{\sigma_n + \sigma_p} E \frac{\partial \Delta n}{\partial x} - \frac{\Delta n}{\tau} = 0 \quad (8.70)$$

Equation (8.70) is the same as the expression (8.43) for the diffusion and drift of non-basic excess charge carriers. Their difference is that there are complex magnitudes that can be defined as D and μ_E , respectively, instead of the diffusion coefficient for the second derivative and the mobility for the first derivative (8.70). $n \approx n_0$, and $p \approx p_0$ Using the Einstein relation for electrons and relativities when at and $\Delta n \ll n_0$ and $\Delta p \ll p_0$ and is $(\mu_n / D_n = \mu_p / D_p = e / kT)$, the ambiquity diffusion coefficient can be written as follows

$$D = \frac{D_n \sigma_p + D_p \sigma_n}{\sigma_n + \sigma_p} = \frac{n_0 + p_0}{\frac{n_0}{D_p} + \frac{p_0}{D_n}} = \frac{kT}{e} \frac{n_0 + p_0}{\frac{n_0}{\mu_p} + \frac{p_0}{\mu_n}}, \quad (8.71)$$

and ambiquity drift mobility can be written as

$$\mu_E = \frac{\mu_n \sigma_p - \mu_p \sigma_n}{\sigma_n + \sigma_p} = \frac{p_0 - n_0}{\frac{n_0}{\mu_p} + \frac{p_0}{\mu_n}}. \quad (8.72)$$

If we use the Einstein relation, the ambipolar diffusion coefficient D can be expressed as follows.

$$D = \frac{kT}{e} \mu_D, \quad (8.73)$$

Comparing this expression with (8.71), the role of mobility is played by magnitude

$$\mu_D = \frac{n_0 + p_0}{\frac{n_0}{\mu_p} + \frac{p_0}{\mu_n}} \quad (8.74)$$

and this is called ambiquity diffusion mobility.

From (8.72) it follows that p_0 and n_0 can change depending on the relationship. This is determined by the description E of the drift velocity of the quasi-neutral cloud of unbalanced charge carriers moving in the direction of motion of non-primary charge carriers in the outer electric field μ_E . In an electronic semiconductor, $\mu_E > 0$ and the cloud moves in the direction of motion of the spurs, while in a p-type material, the cloud moves in the opposite direction in the direction of the movement of electrons.

For a private semiconductor where $n_0 = p_0 = n_i$ we get:

$$D = 2 \frac{D_n D_p}{D_n + D_p} = 2 \frac{kT}{e} \frac{\mu_n \mu_p}{\mu_n + \mu_p}; \quad (8.75)$$

$$\mu_D = 2 \frac{\mu_n \mu_p}{\mu_n + \mu_p} \quad (8.76)$$

$$\mu_E = 0 \quad (8.77)$$

From (8.75) and (8.77) it follows that the value of the ambivalent diffusion coefficient for a private semiconductor depends solely on the diffusion coefficients (or mobility) of electrons and swells. The equality of the ambipolar drift mobility to zero indicates that the external electric field in a private semiconductor, when $\Delta n = \Delta p$, does not affect the spatial distribution of the charge carriers.

For semiconductors with input ($n_0 \gg p_0$ yoki $p_0 \gg n_0$) where the concentration of the main charge carriers significantly exceeds the concentration of non-basic charge carriers, the ambiquity diffusion coefficient D , the ambiquity diffusion mobility μ_D , and the ambiquity drift mobility μ_E It is equal to the diffusion coefficient and mobility of non-basic charge carriers, respectively. For example, for n-type semiconductor ($n_0 \gg p_0$)

$$\begin{aligned} D &= D_p, \\ \mu_D &= \mu_E = \mu_p. \end{aligned} \quad (8.78)$$

Consequently, the diffusion and drift of excess charge carriers in an inlet semiconductor are determined by the diffusion coefficient and the mobility of the non-primary charge carriers, respectively.

Control Questions

1. What do you mean by diffusion of unbalanced charge carriers?
2. What is the drift of unbalanced charge carriers?
3. Explain the diffusion and drift of non-primary excess charge carriers in input semiconductors
4. Give an insight into the diffusion and drift of unbalanced charge carriers in a semiconductor near special conductivity
5. Induce the continuity equation
6. Give an understanding of diffusion and diffusion
7. Cite the Einstein relationship

Chapter IX.

CONTACT EVENTS IN SEMICONDUCTORS

9.1. Semiconductor in External Electric Field

Contact phenomena are based on physical processes that occur in a semiconductor contact under the influence of an electric field. To understand these phenomena, one needs to know the properties of a semiconductor located in the external electric field. For this purpose, we consider the properties of an electron semiconductor incorporated into the homogeneous electric field of a capacitor (Fig. 9.1).

If a semiconductor has no external electric field, the volumetric charge is zero. In the presence of an external electric field, a redistribution of charge carriers occurs in the semiconductor, resulting in a volumetric charge with density $\rho(r)$ and an electric field with a voltage $E(r)$. The change in the distribution of the concentration of charge carriers resulting in the appearance of a volumetric charge occurs in the region adjacent to the surface of the semiconductor. The volumetric charge shields the outer electric field, due to which it penetrates only the surface layer of the semiconductor.

As shown in Figure 9.1, when an external electric field is imposed, the concentration of electrons increases in the area near the surface of the semiconductor (Figure 9.2- a) and the result is a negative volumetric charge (Figure 9.2- b) that the sample penetrated from the surface to the inside

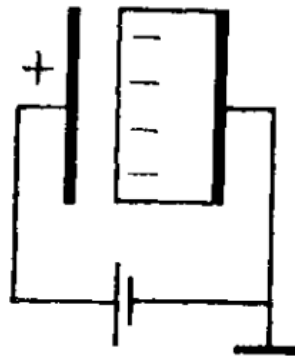


Fig. 9.1. Semiconductor in a homogenous electric field

With increasing distance, it decreases. The negative volume generates a charge electric field, the voltage of which is maximum on the surface of the semiconductor E_s (Fig. 9.2- v). This electric field changes the potential energy of the electron (Fig. 9.2, d) by an amount equal to $U(r) = -e\phi(r)$, wherein is the electrostatic potential of the contact field $\phi(r)$ (Fig. 9.2-g). Thus, the electric field

causes the semiconductor to bend the energy band.

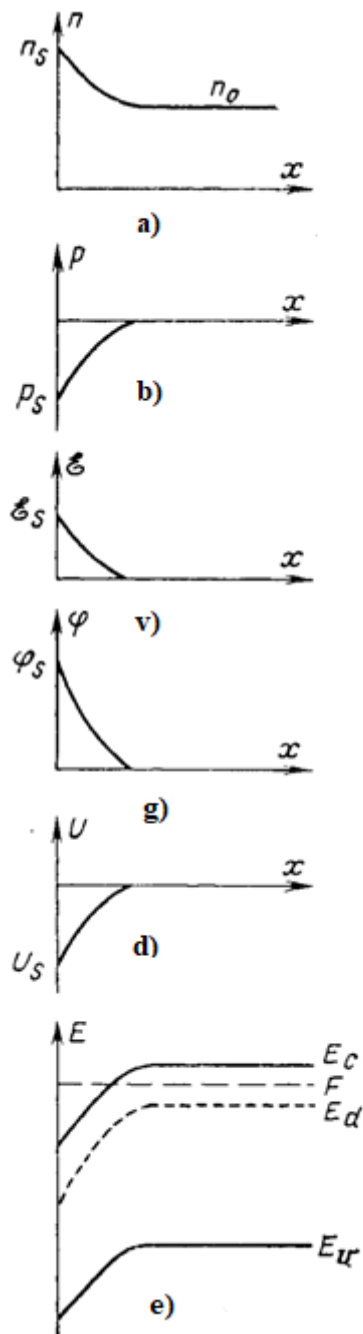


Fig. 9.2. Variations in the concentration of free electrons (a), volumetric charge (b), field strength (v), field potential (g), potential energy of electrons (d) and the curvature of the donor semiconductor zones (e) in the outer electric field.

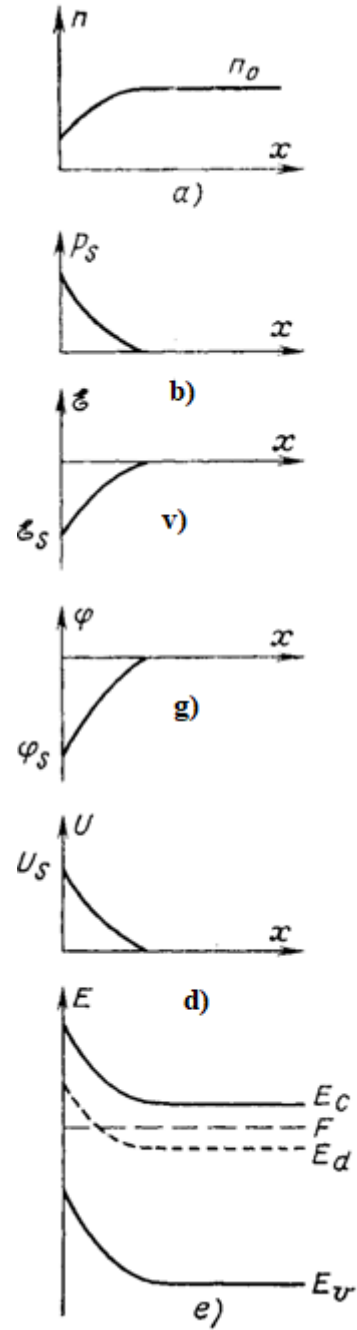


Fig. 9.3. Changes in the concentration of free electrons (a), volumetric charge (b), field strength (v), field potential (g), potential energy of electrons (d) and curvature of donor semiconductor zones (e) when the external electric field direction is changed.

$$\left. \begin{aligned} E_C(\vec{r}) &= E_C - U(\vec{r}); \\ E_V(\vec{r}) &= E_V - U(\vec{r}). \end{aligned} \right\} \quad (9.1)$$

In this case, the states of all energy levels change, including access levels lying in the exclusion zone (Fig. 9.2e).

Since the semiconductor is in a thermodynamic equilibrium state, the state of the Fermi level is constant (see § 5.2), so the distance between the Fermi level and the fixed energy zones varies. This distance without a field:

$$\left. \begin{aligned} E_C - F; \\ F - E_V. \end{aligned} \right\} \quad (9.2)$$

If there is an area, it is as follows:

$$\left. \begin{aligned} [E_C - U(\vec{r})] - F; \\ F - [E_V - U(\vec{r})]. \end{aligned} \right\} \quad (9.3)$$

A comparison of (9.3) and (9.2) shows that if the distance between E_C and F is reduced by the value $U(r)$, the distance between F and E_V increases by the same value. F and the change in the distance between the edges of the energy zones should cause a change in the distribution of electrons across the layers. In the case under consideration, as shown in Fig. 9.2-e, away from the surface, the semiconductor has electron conductivity, is not transmitted, and there are electrons on the donor input surface (because the Fermi level is higher than the input level), and in the area close to the surface, the semiconductor becomes electron, and the same because the Fermi surface is now in the conduction zone.

When the direction of the applied electric field changes in the near-surface area of the electron semiconductor, the concentration of the electron decreases relative to the concentration in the volume of the semiconductor (Fig. 9.3-a) and the signal of the volumetric charge (Fig. 9.3-b) and potential energy change (Fig. 9.3, v) becomes positive. In this case, even the type of primary charge carriers changes in the surface area, as shown in Fig. 9.3-e—the semiconductor has changed from electron to hollow, that is, an inverse or inverted layer appears. At a certain distance from the surface of such a semiconductor, a special conductivity or i-layer sphere is formed, and the Fermi surface is located in the middle of the no-go zone. This area of the semiconductor, which is close to the i-layer, where the type of electric conductivity changes, is called the physical p - n transition. This field disappears when the external electric field is removed.

Energy diagram We look at the effect of an external electric field on a one-dimensional unseparated n-type semiconductor, shown in Fig. 9.3. As we know, the electric field strength E and the volumetric charge density ρ are mutually bound by the Poisson equation.

$$\frac{dE}{dx} = \frac{1}{\varepsilon_r \varepsilon_0} \rho(x) \quad (9.4)$$

If the field strength is expressed by the gradient of the potential $E = -\text{grad}\varphi$, the Poisson equation will look like this:

$$\frac{d^2\varphi}{dx^2} = -\frac{1}{\varepsilon_r \varepsilon_0} \rho(x) \quad (9.5)$$

Let the concentration of electrons in the semiconductor *volume* ($x \rightarrow \infty$) be equal to n_0 . Since the semiconductor is not deflected, in an area close to its surface

$$n = N_C e^{-(E_C + U - F)/kT} = n_0 e^{-U/kT} \quad (9.6)$$

In the surface layer, the volumetric charge is determined by the positive ions and free electrons of the donor input.

The donor input is considered to be fully ionized, i.e. $N_d^+ = n_0$. . In this case, the volumetric charge in the surface layer is equal to:

$$\rho = e(N_d^+ - n) = e(n_0 - n) = en_0(1 - e^{-U/kT}) \quad (9.7)$$

We are limited to the small curvature of zones under the influence of an external electric field, i.e., spreading $|U| \ll kT$. $e^{-U/kT}$ the expression over a series and limiting it to the first state of the series, and based on the relation (9.7), we present the following:

$$\rho = \frac{en_0 U}{kT} = -\frac{e^2 n_0}{kT} \varphi \quad (9.8)$$

$l_e^2 = \frac{\varepsilon_r \varepsilon_0 kT}{e^2 n_0}$ We introduce a designation in which (9.5) looks like this:

$$\frac{d^2\varphi}{dx^2} = -\frac{1}{l_e^2} \varphi \quad (9.9)$$

The solution to this equation is as follows:

$$\varphi = A_1 e^{-x/l_e} + A_2 e^{x/l_e} \quad (9.10)$$

Because $x \rightarrow \infty \varphi \rightarrow 0$ at $A_2 = 0$, $x = 0$ is at the point $\varphi = -\varphi_s$ and $A_1 = -\varphi_s$ so the potential in the area near the surface

$$\varphi(x) = -\varphi_s e^{-x/l_e}, \quad (9.11)$$

Field Strength

$$E(x) = -\frac{d\varphi}{dx} = -\frac{\varphi_s}{l_e} e^{-x/l_e} = -E_s e^{-x/l_e}, \quad (9.12)$$

Potential Energy of Electrons

$$U(x) = -e\varphi(x) = e\varphi_s e^{-x/l_e} = U_s e^{-x/l_e}, \quad (9.13)$$

Volumetric charge density at the surface

$$\rho_s = \frac{en_0 U_s}{kT} \quad (9.14)$$

Thus, when a semiconductor is introduced into an electric field, curvature of energy

zones near the surface and a change in the concentration of electrons and spikes occurs. At $U_s > 0$, the zones bend upwards and the density of the hives increases (Fig. 9.3). In this case, the surface layer of the n-type semiconductor is enriched with non-basic charge carriers, while the p-type semiconductor is enriched with basic charge carriers. At $U_s < 0$, the zones bend downwards (Fig. 9.2-e) and the concentration of electrons increases, i.e., enrichment of the main charge carriers occurs in the n-type semiconductor *and* the enrichment of non-basic charge carriers occurs in the p-type.

(9.10) (9.13) l_e Debay screening length. This means that l_e is the distance at which the intensity of the external electric field entering the matter is $e \approx 2.7$ times. For a metal where $n_0 = 10^{22} \text{ sm}^{-3}$ va $\varepsilon_r = 1$ the Debay shielding length is approximately 10^{-8} cm , while in Germany where $n_0 = 10^{14} \text{ sm}^{-3}$ va $\varepsilon_r = 16$, the electric field is within a distance of 4 mkm .

9.2. Thermoelectronic output operation

As you know, it takes some energy to transfer an electron from a solid into a vacuum. Consequently, at the boundary between solids and vacuum there is an energy barrier (barrier) that prevents electrons from leaving the crystal. Therefore, only electrons with enough energy to overcome this obstacle can leave a solid body. Obviously, the higher the temperature, the more electrons will be able to overcome a potential obstacle. The release of electrons from matter due to thermal excitation is called thermoelectronic emission.

We determine the density of the thermoelectronic emission current of electrons from an unfused semiconductor. To do this, it is necessary to count the number of electrons that are able to overcome a potential barrier and leave the semiconductor. We denote the energy of an electron exiting the semiconductor into a vacuum and in a state at rest relative to the sample, by $E_a = E_{vac}$. In this case, energy χ (Fig. 9.4) is needed to transfer an electron from the bottom of the semiconductor conduction zone into a vacuum without giving it any velocity:

$$\chi = E_a - E_C. \quad (9.15)$$

The energy χ is called the electron proximity energy. In numerical terms, it is equivalent to the work required to transfer an electron from the bottom of the conduction zone to a vacuum without giving it kinetic energy.

In order to overcome a right-angle potential obstacle at the height v_x , the kinetic energy of an electron with velocity v_x must not be less than the depth of the potential barrier, namely:

$$\frac{m^* v_x^2}{2} \geq \chi. \quad (9.16)$$

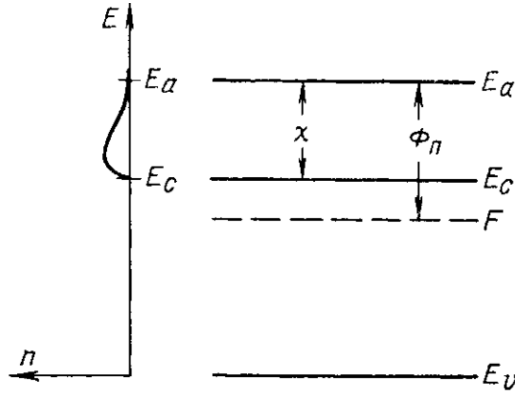


Fig. 9.4. Work of Electron Semiconductor Thermoelectronic Output

The number of quantum states for a unit-size crystal is determined by the following expression in the range of velocities *from* V to $V + dV$ (see § 4.1):

$$dZ = 2\left(\frac{m^*}{h}\right)^3 dV_x dV_y dV_z \quad (9.17)$$

For electrons that can leave the semiconductor, $E - F \gg kT$, their number will be as follows:

$$dn = f_0 dZ = 2\left(\frac{m^*}{h}\right)^3 e^{-(E-F)/kt} dV_x dV_y dV_z \quad (9.18)$$

Let's assume that all the electrons that have overcome the potential barrier do not return to the semiconductor. An electron dnv_x equal to the flow on the crystal surface v_x . If it travels with velocity, then the density of the current flowing from the vacuum to the semiconductor is as follows:

$$J = e \int_{v_x \min}^{\infty} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} dn v_x = 2e \left(\frac{m^*}{h}\right)^3 e^{F/kt} \int_{v_x \min}^{\infty} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} e^{-E/kt} v_x dV_x dV_y dV_z \quad (9.19)$$

The Complete Energy of the Electron

$$E = E_c + \frac{m^*}{2} (v_x^2 + v_y^2 + v_z^2),$$

Taking into account (9.19), the equation can be written as follows:

$$J = 2e \left(\frac{m^*}{h}\right)^3 e^{-(E-F)/kt} \int_{v_x \min}^{\infty} e^{-m^* v_x^2 / 2kt} v_x dV_x \int_{-\infty}^{\infty} e^{-m^* v_y^2 / 2kt} dV_y \int_{-\infty}^{\infty} e^{-m^* v_z^2 / 2kt} dV_z. \quad (9.20)$$

To calculate J we use the following relationship:

$$\int_{-\infty}^{\infty} e^{-ay^2} dy = \sqrt{\frac{\pi}{a}}, \quad (9.21)$$

and (9.16) Integration on v_x provides:

$$\int_{v_{x \min}}^{\infty} e^{-m^* v_x^2 / 2kt} v_x dv_x = \frac{kT}{m^*} e^{-m^* v_{x \min}^2 / 2kt} = \frac{kT}{m^*} e^{-\chi / kt}. \quad (9.22)$$

Based on (9.20) – (9.22), we get the expression for the thermoelectron emission current density:

$$J = \frac{4\pi e m^* k^2}{h^3} T^2 e^{-(E_a - F) / kt} = AT^2 e^{-\Phi / kt}, \quad (9.23)$$

here

$$A = \frac{4\pi e m^* k^2}{h^3}; \Phi = \chi + E_c - F = E_a - F. \quad (9.24)$$

According to Energy F (9.24), which is called the thermoelectron output work, it is equal to the difference between the energy of an electron in a vacuum that is in a quiescent state on the surface of a semiconductor sample and the Fermi surface in the semiconductor. As with metals, F for semiconductors is usually equal to a few electron volts.

Since the state of the Fermi state in a semiconductor depends on the temperature, nature and concentration of the input, the work of the thermoelectronic output is also determined by these parameters. We represent the output work for the various semiconductors using (9.24) and the corresponding relations for the Fermi level obtained in §4.6 and 4.7.

The output case F for a private semiconductor is as follows:

$$\Phi_i = \chi + \frac{1}{2} E_g + \frac{kT}{2} \ln \left(\frac{m_n^*}{m_p^*} \right)^{3/2}. \quad (9.25)$$

From the given expression it follows that the output of an electron from a private semiconductor depends on the width of the forbidden zone, the temperature and the ratio of the effective masses of electrons and spikes.

Inputs for weak ionized donor semiconductor

$$\Phi_n = \chi + \frac{E_c - E_d}{2} + \frac{kT}{2} \ln \frac{GN_c}{N_d}, \quad (9.26)$$

Inputs for Strong Ionized Donor Semiconductor

$$\Phi_n = \chi + kT \ln \frac{N_c}{N_d}. \quad (9.27)$$

If the thermoelectronic output of a strongly ionized donor input is determined by the concentration of the input and the temperature of the semiconductor, in case of weak ionization, it also depends on the depth of the level of donor inputs in the forbidden zone and the degree of rotation.

Inputs to weak and strongly ionized actuator semiconductor

$$\Phi_p = \chi + E_g + \frac{E_v - E_a}{2} - \frac{kT}{2} \ln \frac{GN_v}{N_a}, \quad (9.28)$$

$$\Phi_p = \chi + E_g - kT \ln \frac{N_v}{N_a}. \quad (9.29)$$

From the comparison of the formulas (9.26) (9.29), it follows that the output of electrons in a hollow semiconductor is greater than that of an electron semiconductor of the same nature.

9.3. Metal-to-metal contact. The contact potential difference

Suppose there are two types of insulated metals: M_1 and M_2 . Suppose that the output of the metal M_1 is Φ_1 and the Fermi level is greater than that of F_1 , and the M_2 metal is greater than that of Φ_2 and F_2 (Fig. 9.5a). First, let us consider the processes that occur in the contact plane C (Fig. 9.5-b). Metals begin to exchange electrons at the first moment of contact. In $F_1 > F_2$, electrons pass from metal M_1 to metal M_2 , so a positive charge occurs in metal M_1 , and a negative charge occurs in metal M_2 . Consequently, between metals whose contact plane is perpendicular to C, localized electric field is generated in a layer about $10^{-8}cm$ thick.

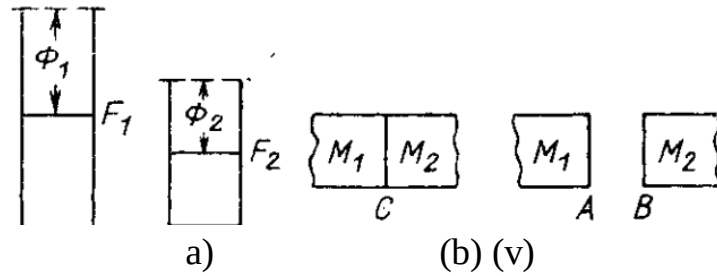


Fig. 9.5. Potential barriers (a) of two insulated metals, contacts of metals in the C-plane (b) and vacuum spaced (v)

The flow of electrons directed from the M_1 metal to the M_2 metal stops when the Fermi level is the same throughout the entire system. From Wikipedia, φ_i the free encyclopedia

To determine φ_i , it is necessary to calculate the number of electrons that have passed from metal M_1 to metal M_2 and vice versa along the x-axis. According to (9.18), electrons with a velocity of v_{x1} are made of M_1 metal.

$$dN_1 = 2 \frac{1}{m_1^* h^3} p_x dp_x \int_{-\infty}^{\infty} \frac{dp_{y1} dp_{z1}}{e^{(E_1 - F_1)/kT} + 1}, \quad (9.30)$$

and the number of electrons passed from the second to the first metal

$$dN_2 = 2 \frac{1}{m_2^* h^3} p_{x2} dp_{x2} \int_{-\infty}^{\infty} \frac{dp_{y2} dp_{z2}}{e^{(E_2 - F_2)/kT} + 1} \quad (9.31)$$

So so.

Since the metal M_1 is positively charged at the C discharge boundary, its electrons receive an additional energy equal to $-e \varphi_i$ (Fig. 9.6-a). But in thermodynamic equilibrium, the total energy of the electrons in the entire system is the same, so

$$E_1 - e\varphi_i = E_2 \quad (9.32)$$

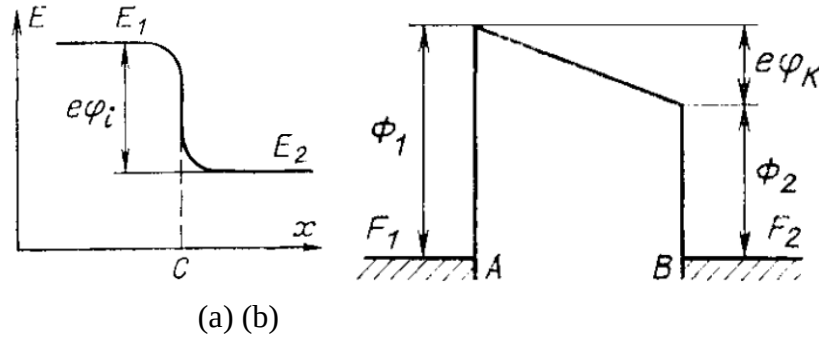


Fig. 9.6. Difference in Internal (a) and External (b) Contact Potentials

Because

$$\frac{\partial E_1}{\partial p_{x_1}} dp_{x_1} = \frac{\partial E_2}{\partial p_{x_2}} dp_{x_2} \quad (9.33)$$

Without it

$$\frac{1}{m_1^*} p_{x_1} dp_{x_1} = \frac{1}{m_2^*} p_{x_2} dp_{x_2} \quad (9.34)$$

Under the conditions of thermodynamic equilibrium, the number of straight transitions is equal to the number of reverse transitions, namely:

$$dN_1 = dN_2. \quad (9.35)$$

This condition can only be fulfilled if the double integrals in (9.30) and (9.31) are equal. They can be equal in the following cases:

$$E_1 - F_1 = E_2 - F_2 = E_1 - e\varphi_i - F_2. \quad (9.36)$$

It follows from this

$$e\varphi_i = F_1 - F_2. \quad (9.37)$$

Thus, the difference in the internal contact potential of two metals in the C plane is determined by the difference in the Fermi energies of the insulated metals.

The Fermi level for metals is associated with the concentration of electrons by the following expression:

$$F_0 = \frac{h^2}{2m^*} \left(\frac{3n_0}{8\pi} \right)^{2/3}, \quad (9.38)$$

So that:

$$e\varphi_i = \frac{h^2}{2} \left(\frac{3}{8\pi}\right)^{2/3} \left(\frac{n_1^{2/3}}{m_1^*} - \frac{n_2^{2/3}}{m_2^*}\right), \quad (9.39)$$

Therefore, the difference in the internal contact potential in the contact of the two metals is determined by the concentration of electrons in the insulated metals and their effective masses. For metals, the value of φ_i is small and is 10^{-3} - 10^{-2} V.

We will now look at the processes that occur in the vacuum friction between metals (AB, Fig. 9.5-v). Since the output of electrons in metal M_1 is greater than that of metal M_2 , at the initial moment of time, the thermoelectron emission current in metal M_2 will be greater than in metal M_1 , i.e. $J_2 > J_1$. This leads to the fact that a positive charge appears in metal M_2 (plane B), a negative charge appears in metal M_1 (plane A), and as a result an electric field appears which prevents the exit of electrons from the metal M_2 . The directed flow of electrons occurs until they reach an equilibrium in the system, which is characterized by a constant Fermi level (Fig. 9.6-b). The difference in contact potentials that arise in a vacuum collision between metals is called the φ_k difference of external contact potentials.

In thermodynamic equilibrium, the equality of emission currents occurs:

$$J_{01} = J_{02} \quad (9.40)$$

or

$$AT^2 e^{-\Phi_1/kT} = AT^2 e^{-(\Phi_2 + e\varphi_k)/kT}, \quad (9.41)$$

In this

$$e\varphi_k = \Phi_1 - \Phi_2. \quad (9.42)$$

Thus, the difference in the external contact potential arising in the vacuum entanglement between metals is expressed by the difference in thermoelectronic output of these metals.

9.4. Metal Semiconductor Contact

Suppose there is a contact between a metal and a semiconductor with an unfused electron. If the Fermi surface of the insulated metal is below the Fermi level of the F_M semiconductor, i.e., $\Phi_M > \Phi_Y$, then at the initial contact moment, the electron flow from the semiconductor exceeds the electron flow from the metal. As a result, the metal in the contact field is negative, the semiconductor is positively charged, and the electric field that arises between the contact samples prevents the transition of E_{bar} electrons from the semiconductor to the metal. The directed flow of electrons continues until the Fermi levels in the system are equal (Fig. 9.7-a), and an equilibrium is established, characterized by the equality of thermoelectron emission currents.

$$J_{M0} = J_{Y0}. \quad (9.43)$$

From this we establish that the difference of contact potentials that arise in

contact of a metal with a semiconductor is equal to:

$$e\varphi_k = \Phi_M - \Phi_Y. \quad (9.44)$$

Electric field strength in the volumetric charge layer

$$E_{bar} = \int \frac{1}{\varepsilon_r \varepsilon_0} \rho dx = \frac{1}{\varepsilon_r \varepsilon_0} Q,$$

Therefore, the difference in potentials in bodies touching each other from the equality of volumetric charges ($Q=\text{const}$) $U = -\int_0^l E_{bar} dx = -\frac{Q}{\varepsilon_r \varepsilon_0} l$ is proportional to the thickness of the volumetric charge layer.

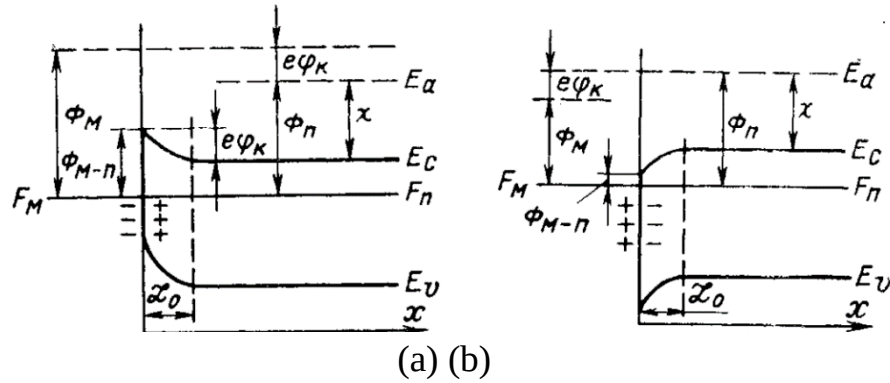


Fig. 9.7. A metal electron semiconductor contact in the cases of $\Phi_M > \Phi_Y$ (a) and $\Phi_M < \Phi_Y$ (b)

But the thickness of the volumetric charge layer in a metal cannot exceed $10^{-8}-10^{-7}$ cm, and in a semiconductor it can be 10^{-4} cm or more, therefore, the difference of potentials in the volumetric charge layer of a metal is very small in comparison with the potential difference in semiconductor. Therefore, the difference in contact potentials φ_k can be considered to be located almost within the contact region of the semiconductor. The electric field strength in the surface layer caused by the contact potential difference does not exceed 10^6 V/cm, and the field strength due to the action of the crystal lattice ions is 10^8 V/cm.

Therefore, a small addition of the contact potential difference cannot change the width of the semiconductor's bandwidth, but its energy band is bent relative to the Fermi level, resulting in a change in the concentration of electrons and spikes in its surface layer.

In this case, if the output of electrons in the semiconductor is less than the output of the electrons in the metal, the semiconductor is positively charged and the energy zone in the contact field bends upwards, so that the lower part of the conductivity zone E_c moves away from the Fermi level F_n and the upper part of the stress zone E_v is the opposite of the Fermi level F_p (Fig. 9.7-A). Therefore, the concentration of electrons near the contact decreases and the number of holes in the

tensile zone increases in relation to their concentration in the semiconductor size. Thus, in the case of $\Phi_M > \Phi_Y$, the concentration of particles in the contact region of an electron semiconductor is greater than its size (Fig. 9.7-a), i.e., a layer with reduced specific conductivity is formed. The contact layer in which the specific conductivity of the semiconductor is reduced (enriched with non-primary charge carriers) is called the closing layer. In the hollow semiconductor, the concentration of the hives in the contact is greater than that of the sample size, i.e. a layer with a higher specific conductivity is formed. The contact layer of a semiconductor with high specific conductivity (enriched with primary charge carriers) is commonly referred to as an antishuter.

If the output performance of a semiconductor is greater than the output work of a metal, i.e. $\Phi_M > \Phi_Y$, the semiconductor is negatively charged and the energy zone in its contact area is bent downwards, so that the number of electrons in the conduction zone near the contact increases and the number of spikes in the stress zone decreases. This means that an antiblocking layer is formed in an electron semiconductor (Fig. 9.7-b), while an antiblocking layer is formed in a semiconductor with an electron.

A physical p-n transition occurs with a strong enrichment of the contact field with non-primary charge carriers, which leads to an inversion of electrical conductivity. Strong enrichment of the contact layer with the main charge carriers can lead to the spinning of the semiconductor in this field.

In a private semiconductor, the curvature of the energy bands in both $\Phi_M > \Phi_Y$ and $\Phi_M < \Phi_Y$ is accompanied by an increase in the specific conductivity in the contact area.

We change (9.44) as follows:

$$e\varphi_k = (\Phi_M - \chi) - (\Phi_Y - \chi) = \Phi_{M-Y} - \Phi_{Y-M}. \quad (9.45)$$

Here $(\Phi_M - \chi) = \Phi_{M-Y}$ is the work that needs to be done to transfer an electron from the Fermi layer of a metal to the lower part of the semiconductor's conductivity zone in contact (see Fig. 9.7); $(\Phi_Y - \chi) = \Phi_{Y-M}$ - The energy that must be spent to transfer an electron in the size of a semiconductor, located at the bottom of the conduction zone, directly into a metal.

Suppose that the electric field enters the electron semiconductor at a depth of L_0 (Fig. 9.7-a). We define it for a semiconductor in which the donor input is fully ionized, i.e. $n_0 = N_d^+$..., since the energy of the electrons at the bottom of the conduction zone in the contact layer $E_c + e\varphi(x)$ is equal to the energy of the electrons, the volumetric charge in this layer according to (9.7) is determined by the relation as follows

$$\rho = en_0[1 - e^{-e\varphi(x)/kT}]. \quad (9.46)$$

Since the difference in contact potentials coincides entirely with the contact field of the semiconductor, $|e\phi(x)| > kT$. The volumetric charge for this case is:

$$\rho = en_0. \quad (9.47)$$

This means that free electrons from a semiconductor field of L_0 length are compressed under the influence of an electric field, and a positive charge is accumulated in the contact field, which is caused by donor input ions. Taking into account (9.47) for the volume charge area, the Poisson equation can be written as follows:

$$\frac{d^2\phi}{dx^2} + \frac{e}{\epsilon_r\epsilon_0}n_0 = 0. \quad (9.48)$$

A common solution to this equation is the following function:

$$\phi(x) = -\frac{en_0}{2\epsilon_r\epsilon_0}(L_0 - x)^2 + A(L_0 - x) + B. \quad (9.49)$$

The field enters the semiconductor only up to the depth of L_0 , and therefore the equation (9.49) must satisfy the following boundary conditions:

$$\begin{aligned} \phi(L_0) &= 0; \\ E(L_0) &= -\frac{d\phi}{dx}\bigg|_{x=L_0} = 0. \end{aligned} \quad (9.50)$$

By putting the marginal conditions in (9.49), we get:

$$\begin{aligned} \phi(L_0) &= B = 0; \\ \frac{d\phi}{dx}\bigg|_{x=L_0} &= -A = 0. \end{aligned} \quad (9.51)$$

Consequently, the electrostatic potential changes with respect to the x coordinate near the contact field of an electron semiconductor (Fig. 9.7-a) as follows:

$$\phi(x) = -\frac{en_0}{2\epsilon_r\epsilon_0}(L_0 - x)^2. \quad (9.52)$$

To determine L_0 , we use the boundary condition of the point $x = 0$. In this state

$$\phi(0) = -\phi_k = -\frac{1}{e}(\Phi_M - \Phi_Y). \quad (9.53)$$

This case allows us to determine the penetration depth of the contact field into the semiconductor based on (9.44) and (9.52):

$$L_0 = \sqrt{2\epsilon_r\epsilon_0\phi_k/en_0} = \sqrt{2\epsilon_r\epsilon_0(\Phi_M - \Phi_Y)/e^2n_0}. \quad (9.54)$$

It follows from this formula that the lower the degree of leation of the semiconductor, the greater the difference in the output of electrons from the metal and the semiconductor, the greater the depth of penetration of the electric field into the semiconductor caused by the difference in contact potentials. From (8.35) and (9.54) $L_0/l_e = \sqrt{2(\Phi_M - \Phi_Y)/kT}$ follows . Therefore, if the output work difference is about 1 eV for a metal semiconductor contact, the penetration depth of the field into

the semiconductor is about 10 times the shielding length for the metal vacuum semiconductor contact.

According to Fig. 9.7-a, the contact layer of a semiconductor, when it is the enclosing layer, the main charge carriers decrease, and therefore the resistance increases relative to the thickness of the semiconductor. Essentially, a metal semiconductor contact has the properties of a capacitor in that it forms a high resistance layer between the semiconductor, which is better conductive than the metal. Therefore, the capacity of the contact perpendicular to the surface unit, which is called the barrier capacity, is equal to:

$$C = \varepsilon_r \varepsilon_0 / L_0 = \sqrt{\varepsilon_r \varepsilon_0 e n_0 / 2 \phi_k}. \quad (9.55)$$

9.5. Current correction on a metal semiconductor contact

Let us return to the consideration of the contact of the metal with the donor semiconductor when the condition $\Phi_M > \Phi_Y$ is met. In the equilibrium state, the bending of the zones in the contact region occurs, resulting in a layer of enclosure in which electrons are reduced (Fig. 9.7-a). Concentration of free carriers at depths in a semiconductor ($x > L_0$)

$$n_0 = N_C e^{-(E_C - F)/kT} = N_C e^{\xi_0/kT}, \quad (9.56)$$

where E_C is the state of the bottom of the semiconductor's conduction zone; $\xi_0 = F - E_C$ is the state of the Fermi layer at the depth of the semiconductor measured from the lower edge of the conduction zone. In thermodynamic equilibrium, the state of the Fermi level is the same for the whole structure:

$$F = E_C + \xi_0 = E_C + e\phi(x) + \xi(x) = \text{const}. \quad (9.57)$$

where $e\phi(x)$ is the potential energy of electrons in the contact region; E , $\xi(x)$ is the state of the Fermi level in the contact field, measured from the bottom of the permeability zone.

When an external voltage is applied to a metal semiconductor contact, the state of the semiconductor becomes unequilibrium and the concentration of conduction electrons is determined by the Fermi quasi-level (see paragraph 7.1) and is equal to:

$$F_k = E_C + e\phi_1(x) + \xi_1(x), \quad (9.58)$$

where $\xi(x)$ and $\phi(x)$ are the quasi-Fermi surface, measured from the bottom of the conduction zone, and the potential energy of the electrons in the contact region of the semiconductor. Therefore, the concentration of free electrons near the contact region of a semiconductor similar to (9.56), i.e. in the field of $0 < x < L$, can be written as follows:

$$n(x) = N_C e^{\xi_1(x)/kT}, \quad (9.59)$$

Since the electron concentration in the contact field *depends on x*, the quasi-Fermi level in the entire field of a semiconductor is not constant. To determine its variation, we calculate the current passing through the semiconductor. Vine Density

$$J = en\mu_n E + eD_n \frac{dn}{dx} \quad (9.60)$$

By means of Fig. 9.59, the current density can be expressed in terms of potential and electron concentrations. As a result, we get:

$$J = n\mu_n \frac{d(e\phi)}{dx} + \frac{eD_n}{kT} \frac{d\xi_1}{dx} \quad (9.61)$$

Given Einstein's attitude, this expression would be:

$$J = n\mu_n \frac{d(e\phi + \xi_1)}{dx}. \quad (9.62)$$

Comparing (9.62) and (9.58), we can write the following:

$$J = n\mu_n \frac{dF_n}{dx}. \quad (9.63)$$

Let's figure out how the quasi-Fermi level changes between two points 1 and 2 where the current densities of a semiconductor are the same:

$$\int_{(1)}^{(2)} dF_n = \int_{(1)}^{(2)} \frac{J dx}{n\mu_n} = eJS \int_{(1)}^{(2)} \frac{dx}{enS\mu_n} = eI \int_{(1)}^{(2)} dR, \quad (9.64)$$

Here $dR = \frac{dx}{\sigma S}$ is the resistance of a semiconductor with a length dx , cross-section C , and a specific electrical conductivity $\sigma = en\mu_n$. Integrating Equation (9.64), we get:

$$\Delta F_n = F_{n_2} - F_{n_1} = eI(R_2 - R_1) = e(U_2 - U_1). \quad (9.65)$$

Thus, when an external electric field is present, the change in the state of the quasi Fermi stahi is determined by the difference in potentials to the deposit.

Now let us determine the density of a current passing through a metal-semiconductor contact for different poles of external voltage. Suppose a negative voltage is applied to a semiconductor and a positive voltage (straight transition) to a metal. Suppose that all external voltage is ordered into the surface layer of a semiconductor which has a high resistance. In this case, the state of the Fermi level in the core of the semiconductor with respect to the E_C level does not change, but with respect to the state of the Fermi level in the metal, it shifts to *the* eU value. Consequently, the difference in contact potentials between a metal and a semiconductor decreases by U and ϕ_k is equal to U , and the state of the quasi-Fermi level near the contact field changes from F_n to F_m . As a result of the reduction of the potential obstacle from the semiconductor side, the equilibrium in the system is disturbed, and a current flows from the metal to the semiconductor and in the opposite direction equal to the difference of thermoelectronic emission currents:

$$J = J_1 - J_2. \quad (9.66)$$

The decrease in the potential barrier on the semiconductor side causes a stronger transition of electrons from the semiconductor to the metal as compared to the equilibrium state. This transition of electrons causes a thermoelectronic emission current directed from metal to semiconductor and is equal to:

$$J_1 = AT^2 e^{-(\Phi_Y + e\phi_k - eU)/kT} = AT^2 e^{-(\Phi_M - eU)/kT}. \quad (9.67)$$

When a positive voltage is applied to a metal, the height of the potential barrier on the metal side does not change, so the flow of electrons from the metal to the semiconductor becomes the same as in the equilibrium state even in the unequilibrium state. *It* generates a thermoelectronic emission current directed from the semiconductor to the metal equal to:

$$J_2 = AT^2 e^{-\Phi_M/kT}. \quad (9.68)$$

The resulting current density passing through the contact is determined by the following expression:

$$J = J_1 - J_2 = AT^2 e^{-\Phi_M/kT} (e^{-eU/kT} - 1) = J_s (e^{-eU/kT} - 1). \quad (9.69)$$

here

$$J_s = AT^2 e^{-\Phi_M/kT} \quad (6.70)$$

The magnitude is called the saturation current density.

If a negative voltage (inverse direction) is applied to the metal, the potential barrier height in the contact region of the semiconductor increases, and the current density expression passing through the contact becomes as follows:

$$J = AT^2 e^{-(\Phi_Y + e\phi_k + eU)/kT} - AT^2 e^{-\Phi_M/kT} = J_s (e^{-eU/kT} - 1). \quad (9.71)$$

Assuming that $U > 0$ for the straight direction and $U < 0$ for the inverse direction, the equations (9.69) and (9.71) can be written as follows:

$$J = J_s (e^{eU/kT} - 1). \quad (9.72)$$

It is concluded from (9.72) that in a direct displacement ($U > 0$) the current density in the contact increases exponentially depending on the voltage, and in the opposite direction ($U < 0$) the current density increases, reaching J_s . Thus, the metal-semiconductor contact has a corrective effect and manifests itself as an alternating current rectifier.

The area of the semiconductor near the contact forms a volumetric charge layer, the thickness of which is expressed by the formula (9.54) in the absence of external voltage. The thickness of the volumetric charge layer is determined when an external voltage is imposed on the contact

$$L = \sqrt{\frac{2\varepsilon_r \varepsilon_0 (\phi_k - U)}{en_0}}, \quad (9.73)$$

Here, as in equation (9.72), the U gesture depends on the pole of the voltage

imposed.

In the right transition ($U > 0$), the thickness of the volumetric charge layer is less than the thickness L_0 of the latency layer in thermodynamic equilibrium $L_{to'g}$, and in the inverse ($U < 0$) direction $L_{tes} > L_0$.

9.6. Diode theory of current correction

The saturated current characteristic of a metal semiconductor contact volt-ampere can be calculated on the basis of diode and diffusion theories.

If the free travel path of the electrons is significantly smaller than the thickness of the barrier layer L_0 , the charge carriers passing through this layer will scatter several times. In this case, the diffusion and drift promoters of the current must also be taken into account. If the electric field is weak, it will be possible to use Ohm law to find the drift current, according to its condition (6.184):

$$E \ll \frac{kT}{el}. \quad (9.74)$$

The maximum field voltage E_s is at the boundary of the semiconductor to the metal ($x = 0$). Taking into account (9.52) and (9.54)

$$E_s = E_{\max} = -\frac{d\varphi}{dx} \Big|_{x=0} = -\frac{en_0L_0}{\varepsilon_r\varepsilon_0} = -\frac{2\varphi_k}{L_0}. \quad (9.75)$$

By putting the value of E_s found at (9.74), we get the condition for the validity of the diffusion theory of correction:

$$L_0 \gg \frac{2e\varphi_k}{\delta kT} l. \quad (9.76)$$

If the ratio of the free travel path of electrons l and the thickness of the enclosing layer L_0 satisfies an inverse inequality,

$$l \gg \frac{\delta kT}{2e\varphi_k} L_0, \quad (9.77)$$

In it, it can be assumed that the charge carriers fly through the covering sphere almost without noticing collisions. In this case, the layer that masks the volumetric charge of the semiconductor resembles the vacuum gap between the electrodes of the electronic lamp. In this connection, the theory used in this case to calculate the volt-ampere characteristic of a metal semiconductor contact is called diode theory.

Diode theory according to (9.77)

$$\frac{\delta kTL_0}{2e\varphi_k l} \ll 1. \quad (9.78)$$

If it does, it does. This condition is fulfilled by the smaller L_0 and the greater the l , the better. But as can be seen from (9.54), the smaller the thickness of the blanket layer, the greater the concentration of free electrons in equilibrium n_0 and the greater the free running path of electrons (6.29), the greater their mobility.

Consequently, the diode theory of current correction in metal-semiconductor contact applies to semiconductors with high concentrations and charge carrier mobility.

For example, for electron-type germanium, $\mu_n = 3900 \text{ cm}^2/(\text{V}\cdot\text{s})$; $\varepsilon_r = 16,5$; $m_n^* = 0,3m_0$ $n_0 = 10^{14} \text{ cm}^{-3}$ is $e\phi_k = 0,3 \text{ eV}$, $L_0 = 2,3 \cdot 10^{-4} \text{ cm}$ and $l = 1,5 \cdot 10^{-5} \text{ cm}$ and the ratio is $kTL_0/2e\phi_k l = 0,6 < 1$. If the electron concentration is increased to $n_0 = 10^{16} \text{ cm}^{-3}$, then $kTL_0/2e\phi_k l = 0,06$. In this case, the use of the diode theory of correction fully justifies itself.

The situation is different in semiconductors, where the concentration and mobility of charge carriers are small. For example, in p-type copper oxide $\mu_p = 60 \text{ cm}^2/(\text{V}\cdot\text{s})$; $\varepsilon_r = 8,75$; $n_0 = 10^{14} \text{ cm}^{-3}$ and $e\phi_k = 0,5 \text{ eV}$, $L_0 = 2,2 \cdot 10^{-4} \text{ cm}$ and $l = 4 \cdot 10^{-7}$ and the ratio is $kTL_0/2e\phi_k l = 15 > 1$. In this case, to calculate the volt-ampere characteristic (VAX) of a metal-semiconductor contact, it is necessary to use the diffusion theory of correction.

In a review of the contact of a metal with a semiconductor (§9.4), it was believed that only electrons with an energy greater than the value Φ_m were able to leave the metal. However, at normal temperatures, the number of electrons with such high energy disappears. But if there is a very thin potential barrier between the samples, the electrons can overcome it as a result of the tunneling effect.

Based on quantum mechanics, the transparency of a potential barrier whose height is Φ_0 and whose thickness is d is expressed by the following formula D.

$$D = e^{-\frac{2d}{h}[2m^*(\Phi_0 - E_x)]^{1/2}}. \quad (9.79)$$

Here $E_x = \frac{m^* v_x^2}{2}$; v_x is the velocity of electrons in the direction of the sample surface. $\Phi_0 \approx 1 \text{ eV}$; Supposing that; From the condition $d = 10^{-8} \text{ cm}$, $E_x \ll \Phi_0$, we get $D = 1/3$. This means that such thin barriers are almost transparent to electrons. This means that in the future, we assume that all electrons whose energy exceeds the value of Φ_{M-Y} are capable of transitioning to a semiconductor.

To move from semiconductor to metal, e electrons ($\phi_k - U$) must have enough velocity in the direction x to overcome the potential obstacle in height. The current generated by such electrons, according to (9.23), is expressed as:

$$J_1 = e \int_{v_x \min}^{\infty} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} dv_x = \frac{4\pi m^* (kT)^{3/2}}{h^3} e^{(F-E_C)/kt} e^{-e(\phi_k - U)/kt}. \quad (9.80)$$

Here is the expression of the concentration of electrons in a thin semiconductor in equilibrium

$$n_0 = 2(2\pi m kT / h^2)^{3/2} e^{(F-E_C)/kt} \quad (9.81)$$

and their average heat velocity

$$v_0 = (8kT / \pi m^*)^{1/2}, \quad (9.82)$$

If we put it, we get:

$$J_1 = \frac{1}{4} en_0 v_0 e^{-e(\phi_k - U)/kt} = \frac{1}{4} en_s v_0 e^{eU/kt}. \quad (9.83)$$

here

$$n_s = n_0 e^{-e\phi_k/kt} \quad (9.84)$$

U = the concentration of electrons on the surface of a semiconductor at 0.

On the metal side, the height of the potential barrier remains the same, so the flux of electrons from metal to semiconductor does not depend on the put voltage, and is equal to the flux of electrons from the semiconductor in the thermodynamic equilibrium state.

$$J_2 = J_1|_{U=0} = \frac{1}{4} en_s v_0. \quad (9.85)$$

The density of the contact-passing aggregate currents is determined by the following expression

$$J = J_1 - J_2 = \frac{1}{4} en_s v_0 (e^{eU/kt} - 1) = J_s (e^{eU/kt} - 1), \quad (9.86)$$

here

$$J_s = \frac{1}{4} en_s v_0. \quad (9.87)$$

This expression is saturated current in diode theory of rectification, which, based on (9.87), is determined by the thermal velocity of electrons and their concentration on the surface of the semiconductor.

9.7. Diffusion theory of current straightening

In the diffusion theory of straightening, both drift and diffusion constituents of the current are taken into account when determining the current passing through the metal-semiconductor contact. In this case, to calculate the volt-amp characteristic (VAX), we proceed as follows. Multiply both parts of (9.60). $e^{e\phi(x)/kt}$ Using Einstein's relativity and potentially expressing field strengths, we get:

$$J e^{e\phi(x)/kt} = n \mu_n e^{e\phi(x)/kt} \frac{d(e\phi)}{dx} + \mu_n k T e^{e\phi(x)/kt} \frac{dn}{dx} = \mu_n k T \frac{d}{dx} [n e^{e\phi(x)/kt}]. \quad (9.88)$$

We integrate the two parts of (9.88) in the range 0 to L with respect to x:

$$J \int_0^L e^{e\phi(x)/kt} dx = \mu_n k T \int_0^L \frac{d}{dx} [n e^{e\phi(x)/kt}] dx. \quad (9.89)$$

From here we express the current density J, which is independent of x:

$$J = \frac{\mu_n k T [n(L) e^{e\phi(L)/kt} - n(0) e^{e\phi(0)/kt}]}{\int_0^L e^{e\phi(x)/kt} dx}. \quad (9.90)$$

We use boundary conditions to calculate (9.90). The size of the semiconductor doesn't have field, so

$$\varphi(L) = 0; \quad n(L) = n_0, \quad (9.91)$$

At the boundary of separation from the metal ($x=0$)

$$\varphi(0) = \varphi_k - U. \quad (9.92)$$

Consider the case where the external voltage is relatively small and there is always a covering layer in the contact area of the semiconductor:

$$(\varphi_k - U) \gg kT / e. \quad (9.93)$$

Under these conditions, the concentration of electrons on the surface of the semiconductor is almost not dependent on the voltage, so we would not be making a big mistake if:

$$n(0) = n_s = n_0 e^{-e\varphi_k / kT}. \quad (9.94)$$

Taking into account the boundary conditions (9.91), (9.92), and (9.94), the expression for the current density (9.90) takes the following form:

$$J = \frac{n_0 \mu_n kT (1 - e^{-eU / kT})}{\int_0^L e^{e\varphi(x) / kT} dx}. \quad (9.95)$$

To calculate the value of the fraction, we replace the integral variable (9.95):

$$\int_0^L e^{e\varphi(x) / kT} dx = \int_{\varphi_k - U}^0 e^{e\varphi(x) / kT} \left(\frac{d\varphi}{dx}\right)^{-1} d\varphi. \quad (9.96)$$

The greatest contribution to the integral is made by the sphere where the value of $\varphi(x)$ is maximum, i.e., by the sphere close to $x = 0$. As a result of the decrease in $\varphi(x)$ with the increase of x , the function under the gesture of the integral decreases sharply, and its contribution can be ignored. Therefore, to compute the integral equally, we can exclude the value at point $x = 0$ $(d\varphi/dx)^{-1}$ outside the integral sign:

$$\left(\frac{d\varphi}{dx}\right)^{-1} = -E_s^{-1}. \quad (9.97)$$

Based on (9.97), we get:

$$\int_{\varphi_k - U}^0 e^{e\varphi(x) / kT} \left(\frac{d\varphi}{dx}\right)^{-1} d\varphi = -E_s^{-1} \frac{kT}{e} [1 - e^{e(\varphi_k - U) / kT}] = E_s^{-1} \frac{kT}{e} e^{e(\varphi_k - U) / kT}. \quad (9.98)$$

We have ignored the unit in parentheses here, since (9.93) is conditional

$$e^{e(\varphi_k - U) / kT} \gg 1. \quad (9.99)$$

Putting the expression (9.98) in (9.95), we find:

$$J = en_0 \mu_n E_s e^{-e\varphi_k / kT} (e^{eU / kT} - 1). \quad (9.100)$$

Using (9.94), for the diffusion theory of the correction, we obtain the volt-ampere characteristic expression of a metal semiconductor contact:

$$J = en_s \mu_n E_s (e^{eU / kT} - 1) = J_s (e^{eU / kT} - 1). \quad (9.101)$$

Here, in contrast to the saturated current density in diode theory, the magnitude that determines the current value in large inverse directions $J_s = en_s \mu_n E_s = en_s v_s$ depends on the put voltage, that is, it is expressed by the drift velocity of the electron on the surface of the semiconductor.

The VAXi properties of the metal-semiconductor contact for diode and diffusion theories are shown in Fig. 9.8 (curves 1 and 2, respectively).

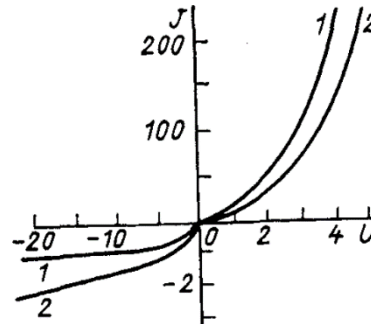


Fig. 9.8. Volt-ampere characteristics of a metal-semiconductor contact:
1st diode theory; 2-Diffusion theory (in relative units)

9.8. Contact of Electron and Hollow Semiconductors

Let us consider the phenomena that occur in the contact of electron and hollow semiconductors. To create such a contact, the donor and acceptor inputs are inserted into the semiconductor. In this case, the concentration of donor and acceptor inputs changes in such a way that in one part of the sample there are donors and acquire electron conductivity, while in the other part containing the receptors and acquire hiccup permeability, as a result of which the conductivity in a certain region of the crystal changes from electron to spike. Such a transition between materials with n- and p-type electrical conductivity is called the p-n transition.

Suppose that the p-n transition is infinitely narrow and the acceptor field of the semiconductor is more strongly legged than the electron field, i.e., $N_a > N_d$ (Fig. 9.9-a).

Let the main charge carriers in the p-field be the concentration of the hives p_p and the non-basic charge carriers - let the electrons be n_p and the concentration of electrons in the n-field be n_n and the concentration of electrons in the n-field be n_n and the density of the hives p_n . Also, let us assume that the energy levels of the donors and the actors are located so close to the conductivity and stress zones that they are all fully ionized. In this case, the following expression is appropriate for the state of thermodynamic equilibrium when $p_p = N_a$, $n_n = N_d$, and uneven:

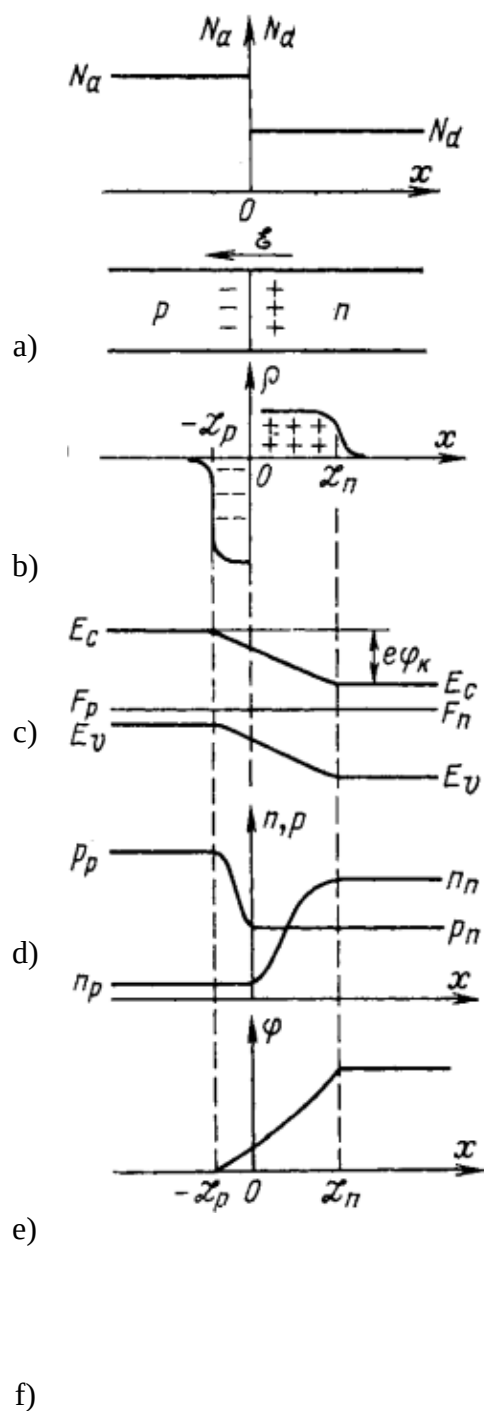


Fig. 9.9. In the contact of electron and hollow semiconductors, the distribution of inputs (a), the dissociation of charge and the appearance of an electric field (b), the distribution of volumetric charge (c), the structure of the zones (d), the distribution of electrons and cavities concentrations (e), and the potential variation (f).

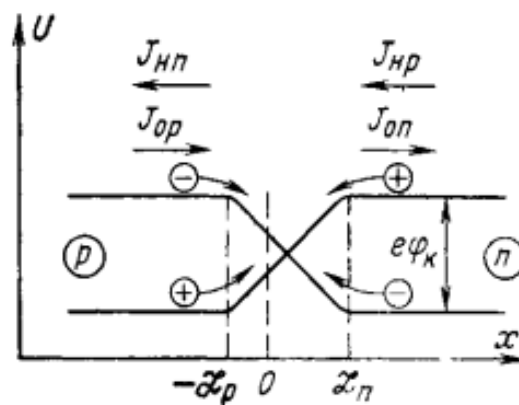


Fig. 9.10. Potential barriers to pitfalls and electrons in the P-n transition:

$$p_p n_p = n_n p_n = n_i^2. \quad (9.102)$$

At the first moment of contact between the n - and p -spheres, a large gradient of electron and p -concentration is formed near the transition boundary. As a result, the diffusion of electrons from the n -sphere to the p -sphere and the holes from the p -sphere to the n -sphere occurs. The appearance of diffusion fluxes leads to the dissociation of the charge, as a result of which a positive volumetric charge is

generated due to the positive ions of the positive donor input in the n -region adjacent to the p - n transition and a negative charge due to negative ions of the actuator input in the p -region near the p - n transition (Fig. 9.9-b, c of the p - n transition). These volumetric charges in the contact field create a strong electric field from the n -field to the p -field and impede the movement of electrons and spurs. As a result, an equilibrium state, characterized by the continuity of the Fermi level for the entire semiconductor, is established, and energy zones are curved in the transition region where there is an electric field (Fig. 9.9-d). The curvature of the energy zones leads to a redistribution of electron and p - n concentrations (Fig. 9.9-e) and changes the direction of the electrostatic potential in the p - n transition region (Fig. 9.9-f).

As can be seen from Fig. 9.10, the main charge carriers must overcome a potential obstacle of $e\varphi_k$ altitude when passing through contact. The transition of non-primary charge carriers occurs under the action of the electric field of the p - n transition. In the case of thermodynamic equilibrium, the densities of the diffusion currents of the main charge carriers $J_{asos,n}$, and $J_{asos,p}$. The densities of the drift currents of the non-primary charge carriers are equal to $J_{asosb,n}$ and $J_{asosb,p}$, and the sum current passing through the p - n transition is zero.

Suppose that Φ_n is the thermoelectronic output of the semiconductor and Φ_p is the thermoelectronic output of the semiconductor. In it, the potential barrier in p - n transition in thermodynamic equilibrium is determined by the following equation:

$$e\varphi_k = \Phi_p - \Phi_n = (\chi - E_C + F_p) - (\chi - E_C + F_n) = F_p - F_n. \quad (9.103)$$

Since the donor and actor inputs are fully ionized, using (9.27) and (9.29) or (4.127) and (4.136) we get the following:

$$e\varphi_k = E_g - kT \ln \frac{N_V}{N_a} - kT \ln \frac{N_C}{N_d}. \quad (9.104)$$

Follow

$$n_i^2 = N_C N_V e^{-E_g/kT}, \quad (9.105)$$

Taking into account (9.104), we get the following:

$$e\varphi_k = kT \ln \frac{n_n p_n}{n_i^2} \quad (9.106)$$

or

$$\frac{p_n}{p_p} = \frac{n_p}{n_n} = e^{-e\varphi_k/kT}. \quad (9.107)$$

Thus, the difference in contact potentials in the p - n transition is greater the greater the n and p fields of the semiconductor are strongly alloyed; its maximum value for an unseparated semiconductor is as follows:

$$\varphi_k = \frac{1}{e} E_g. \quad (9.108)$$

Based on (9.107), the following relationships can be obtained for the concentration of primary and non-primary charge carriers:

$$p_n = p_p e^{-e\varphi_k / kT}; n_p = n_n e^{-e\varphi_k / kT}. \quad (9.109)$$

The output of the electrons in the semiconductor is very large (a few electron volts), and there are almost no electrons having enough energy to leave the crystal at room temperature. At the same time, electrons can go directly from the n-field to the p-sphere, overcoming a potential obstacle. For example, for germanium, the values of $T=300\text{ K}$ and concentration values are $N_d=n_n=10^{16}\text{ cm}^{-3}$; $p_n = 10^{10}\text{ cm}^{-3}$; $N_a = p_p = 10^{14}\text{ cm}^{-3}$; $n_p = 10^{12}\text{ cm}^{-3}$ the potential barrier is:

$$e\varphi_k = kT \ln \frac{p_p}{p_n} = 0,25\text{ eV} \quad (9.110)$$

Since the p-sphere is more strongly charged than the n-sphere, $p_p \gg n_n$, and therefore the thickness of the volumetric charge layer in the small Omli field L_p is less L_n than in the L_p electron sphere, i.e., $L_n > |L_p|$ (Fig. 9.12- c) and the entire volumetric charge areas are as follows:

$$L_0 = L_n + L_p. \quad (9.110)$$

$-L_p \leq x < 0$ In the interval, the volumetric charge is negative and is determined by the concentration of actuator input ions:

$$\rho = -eN_a^- = -ep_p \quad (9.111)$$

and for this area, the Poisson equation looks like this:

$$\frac{d^2\varphi}{dx^2} = \frac{ep_p}{\varepsilon_r \varepsilon_0}. \quad (9.112)$$

$0 < x < L_n$ The volume charge in the range is positive and is determined by the ions of donor inputs:

$$\rho = eN_d^+ = en_n \quad (9.113)$$

and for this area, the Poisson equation looks like this:

$$\frac{d^2\varphi}{dx^2} = -\frac{en_n}{\varepsilon_r \varepsilon_0}. \quad (9.114)$$

In this case, the limit conditions will be as follows:

$$\varphi(-L_p) = 0; \frac{d\varphi}{dx} \Big|_{x=-L_p} = 0; \quad (9.115)$$

$$\varphi(L_n) = \varphi_k; \frac{d\varphi}{dx} \Big|_{x=L_n} = 0; \quad (9.116)$$

Equations (9.112) and (9.114) are the solution according to the terms (9.52) F For the following expressions are satisfied: at $-L_p \leq x < 0$

$$\varphi_p = \frac{ep_p}{2\varepsilon_r\varepsilon_0}(L_p + x)^2; \quad (9.117)$$

At $0 < x < L_n$

$$\varphi_n = \varphi_k - \frac{en_n}{2\varepsilon_r\varepsilon_0}(L_n - x)^2. \quad (9.118)$$

Potential at $x = 0$ (Fig. 9.12-e) and its derivative is continuous, so that

$$\varphi_p(0) = \varphi_n(0); \quad (9.119)$$

$$\left. \frac{d\varphi_p}{dx} \right|_{x=0} = \left. \frac{d\varphi_n}{dx} \right|_{x=0} \quad (9.120)$$

By applying the condition (9.120) to expressions (9.117) and (9.118), we get:

$$n_n L_n = p_p L_p. \quad (9.121)$$

Consequently, the volumetric charges in both fields adjacent to the p - n pass of the semiconductor are equal. This is a condition of maintaining electrical neutrality.

It is easy to get the following relations from (9.121):

$$\frac{L_n}{L_0} = \frac{p_p}{n_n + p_p}; \quad \frac{L_p}{L_0} = \frac{n_n}{n_n + p_n}. \quad (9.122)$$

Using these terms, based on equations (9.117)-(9.119), we obtain:

$$\varphi_k = \frac{e}{2\varepsilon_r\varepsilon_0}(n_n L_n^2 + p_p L_p^2) = \frac{e}{2\varepsilon_r\varepsilon_0} L_0^2 \frac{n_n p_p}{n_n + p_p}. \quad (9.123)$$

From this we find the expression of the total thickness of the volumetric charge layer of the p - n transition:

$$L_0 = \sqrt{\frac{2\varepsilon_r\varepsilon_0}{e} \varphi_k \frac{n_n + p_p}{n_n p_p}}. \quad (9.124)$$

As can be seen from this formula, the higher the degree of lequefaction in the n - and p -fields of the semiconductor, the smaller the thickness of the L_0 volumetric charge region. If one of the fields is significantly more strongly depleted than the other, then most of the decrease in the electrostatic potential occurs in the high-resistance field (Fig. 9.9-f).

As can be seen from Fig. 9.9-d, the concentration of charge carriers in the p - n transition region decreases significantly. Therefore, the resistance of the p - n transition is higher when compared to the resistance of an n - or p -type semiconductor layer of the same surface and thickness in the volumetric charge region.

Therefore, the electron-cavity transition is a low-conductivity layer located between high conductivity fields, and has capacitor properties. The capacity per unit of surface, called barrier capacity, can be defined by the formula:

$$C = \frac{\varepsilon_r \varepsilon_0}{L_0} = \sqrt{\frac{e \varepsilon_r \varepsilon_0}{2 \varphi_k} \frac{n_n p_p}{n_n + p_p}}. \quad (9.125)$$

9.9. P-N Equalization at Transition

Consider a P-n transition where the p-domain is positively charged (straight transition) (Fig. 9.11- b) with a potential difference. Since the resistance of the p-n transition volumetric charge layer is high, the voltage drop is mainly in this area. Therefore, with a straight transition, the height of the potential barrier decreases by eU relative to the equilibrium state (Fig. 9.11-a), and the thickness of the obstacle layer changes accordingly. To calculate it, you need to use the following expression instead of the formula (9.124):

$$L_{to'g'} = \sqrt{\frac{2 \varepsilon_r \varepsilon_0 (\varphi_k - U)}{e} \frac{n_n + p_p}{n_n p_p}}. \quad (9.126)$$

A decrease in the potential barrier leads to an increase in the flow of primary charge carriers relative to the equilibrium state, since a large number of electrons from the n-sphere pass to the p-sphere and a large number of electrons from the p-field to the n-sphere, but the flow of non-primary charge carriers is almost unchanged

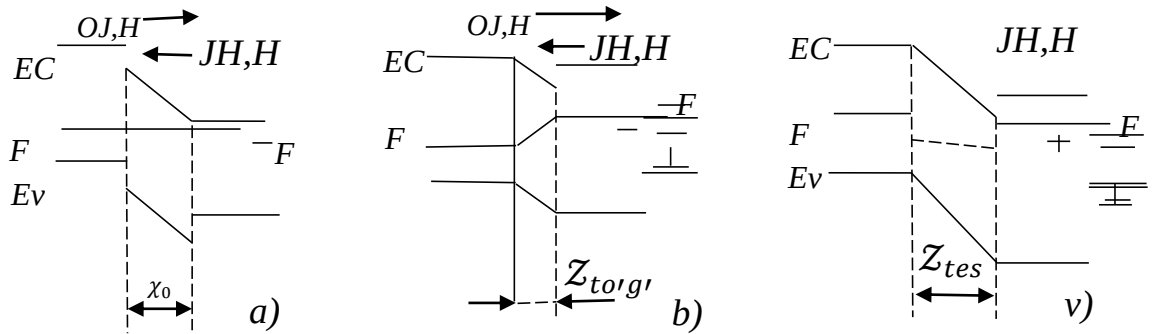


Fig. 9.11. Energy diagram of the p-n transition in thermodynamic equilibrium is (a), the right (b), and the inverse (v) transitions

As a result, in the outer circuit, the difference in the currents of the primary and non-primary charge carriers is equal to the difference and a current ("straight current") flows from the p-sphere to the n-sphere (Fig. 9.11-b), i.e.

$$J = J_{asos} - J_{asoiyemas}. \quad (9.127)$$

The redundant primary charge carriers that appear in the n-field — the spikes Δp — generate a positive volumetric charge near the contact at the first minute. However, after a very short period of time, as determined by Maxwell's relaxation time (see paragraph 7.3), this charge is compensated by the volumetric charge of the main charge carriers—electrons, which rise by n from the depth of the Δn -field

under the influence of the electric field created by the excess spurs , and the electrons rise by n -enters the industry from an external chain. Thus, electric neutrality is maintained in all parts of the electron semiconductor, but the concentration of electrons and spikes in the contact region of the p - n transition increases by $\Delta n = \Delta p$ with respect to the equilibrium state. The introduction of non-core charge carriers into the semiconductor field using the application of a straight p - n transition is called an injection. Now the concentration of hives in the n -sphere near the contact will be as follows:

$$p = p_n + \Delta p. \quad (9.128)$$

To find it in the stationary state $x=L_n$ and the unequilibrium state of the charge carriers, one must use the value $e(\varphi_k - U)$ instead of $e\varphi_k$ in (9.109). So,

$$p(L_n) = p_p e^{-e(\varphi_k - U)/kT} = p_n e^{eU/kT}, \quad (9.129)$$

It follows that the concentration of excess hives in the n -area at $x=L_n$ is:

$$\Delta p(L_n) = p - p_p = p_n (e^{eU/kT} - 1). \quad (9.130)$$

Similar phenomena occur in the p -domain: electrons are injected here from the n -sphere, and the concentration of excess electrons in $x=-L_p$ is as follows:

$$\Delta n(-L_p) = n_p (e^{eU/kT} - 1). \quad (9.131)$$

It follows from (9.130) and (9.131) that with the increase in the right pass in the P - n transition, the concentration of injected non-core charge carriers increases dramatically, resulting in a strong increase in the current passing through the contact in the right direction.

If an inverse transition is applied to the p - n transition (Fig. 9.11-c), the p -sphere is negatively charged, the potential barrier increases by the amount of eU , and the thickness of the volumetric charge masking layer increases. L_{tes} is calculated by putting the value $\varphi_k + U$ in accordance with (9.124). The more the transition moves in the opposite direction, the smaller the number of primary charge carriers capable of overcoming the growing potential obstacle. Accordingly, with respect to the equilibrium state, the number of non-primary charge carriers in the contact area decreases, and therefore the number of most charge carriers also decreases due to the observance of electric neutrality. This phenomenon is termed charge carrier extraction. The concentration of excess electrons in the p -sphere is still determined by equation (9.131), but the U -value is now taken by a negative hint. Thus, in the reverse transition of the p - n transition, the current of the main charge carriers is less than the equilibrium state, and the current of the non-primary charge carriers is almost unchanged. Therefore, by passing p - n , the total current is directed from the n -sphere to the p -field, and with the increase in the reverse voltage increases a little at first, and then tends to a certain value, which is called the saturation current. Consequently, the p - n transition has a nonlinear VAX (Fig. 9.12).

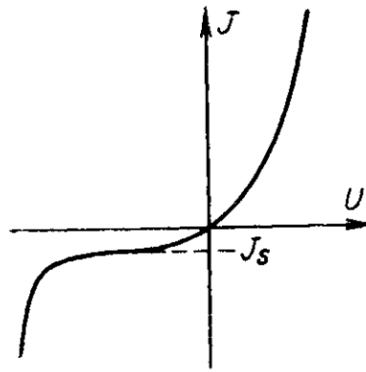


Fig. 9.12. Volt-ampere characteristic of a p-n transition

Control Questions

1. What can be said about semiconductor in outdoor electric field?
2. What do you mean by a thermoelectron output operation?
3. What is thermoelectron emission?
4. What processes take place in metal-to-metal contact?
5. Explain the difference in contact potentials
6. Give an Insight on the Metal-Semiconductor Contact
7. What do you mean by straightening?
8. Describe the diode theory of current correction
9. Give an Insight into the Diffusion Theory of Current Straightening
10. Describe the processes that take place in the contact of semiconductors with electrons and cavities
11. How is the current straightened in the P-n transition?

Chapter X

SURFACE PHENOMENA IN SEMICONDUCTORS

10.1. The nature of surface layers

The emergence of interest in semiconductor surface physics has been the fact that the effects of surface phenomena on processes in the semiconductor volume, as well as on the operation of most semiconductor instruments, have been identified. Currently, the large-scale production of structures in semiconductor electronics consisting of integrated circuits containing a large number of transistors, capacitances, inductances and other elements, their introduction in various fields (computers, TVs, communications and power devices) makes the issue of a serious study of the surface physics of semiconductors urgent, due to the ever-shrinking size of the semiconductor elements used. The fact that they are increasing the importance of the surface day by day.

A semiconductor crystal has a system of ring (discrete) and continuously distributed energy layers on its surface. The reasons for their occurrence vary, below we will look at some of them.

Local energy levels on the surface and the charge in them occupy energy fields near the surface. And this makes the crystals change electrical conductivity, output work, photoelectric properties.

Studying the surface effects of semiconductors can help determine most phenomena. The state of the surface indicates how the processes that take place at the crystal size take place. For example, surface surfaces are centers of recombination, and unbalanced charge carriers in the sample may affect the residence time.

In 1932, I.E. Tamm showed that the energy spectrum of electrons in an infinite and finite crystal differs significantly. From the fields of allowed energies of electrons in an infinite crystal, new energetic states for electrons arise as a result of its disconnection in a finite crystal. The wave functions corresponding to these states are the greatest (maximum) at the crystal boundary, decaying as they travel into the crystal and toward the vacuum side.

If the crystal is bounded and the potential barrier in the edge elementary cell is different from that inside the crystal, then permissible energy states can occur on the finite crystal surface instead of the forbidden states within the infinite crystal. As we know, in order to generate the energy spectrum of any quantum system, the Schrödinger equation must be solved. To do this, it is enough to know the potential change within the crystal in addition to the choice of wave function (Fig. 10.1).

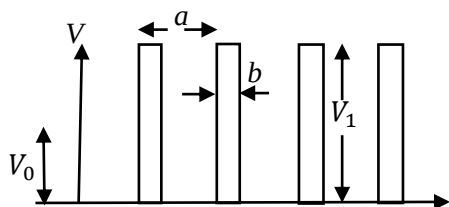


Fig. 10.1. The Kronig-Penny model of weakly bonded electrons

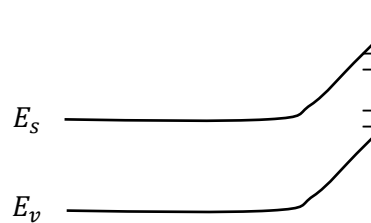


Fig. 10.2. Appearance of energy fields on the surface

E.I. Tamm in his first account adopted the Kronig-Penny model of weakly bound electrons in solving a one-dimensional problem. At the same time, there was a difference between the marginal potential and the internal barriers. In this way, Tamm created an equation for determining the energy of the layers. In the subsequent development of the theory of surface electron states, the theory generalized and studied the state of three-dimensional space, as well as various manifestations of the periodic potential within a crystal.

As a result of the interruption of the periodic arrangement of crystal atoms on a finite regular crystal surface, a spectrum of ring-ring electron states is formed. These conditions are called Tamm cases (levels).

The density (concentration) of the tamm states is determined using the density of atoms on the surface, i.e., this density is on the order of 10^{15} cm^{-2} . When such density is present, energetic fields on the surface appear (Fig. 10.2).

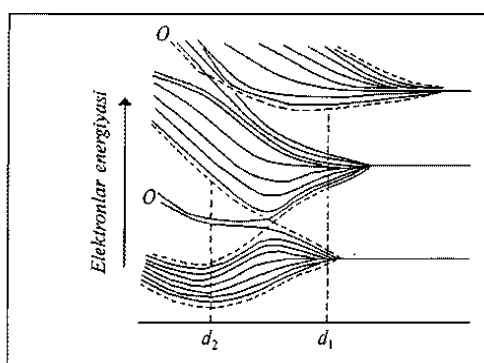


Fig. 10.3. Formation of zones in the crystal

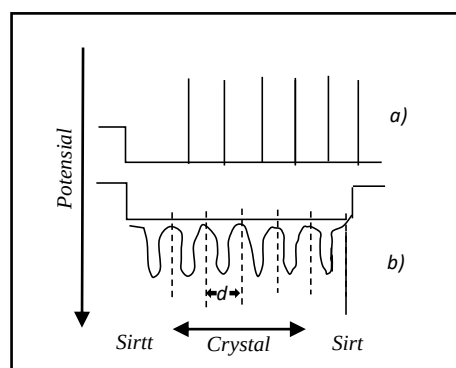


Fig. 10.4. Appearance of two levels in the forbidden area of the crystal

Shock-type surface layers, like Tamm layers, occur as a result of the limitation of the periodic potential on the surface. But the reasons for the origin of the Tamm and Shocky surface layers are different. The shock model looked at a one-dimensional finite chain of 8 atoms. The periodic potential is imagined in the form

of voluntary potential pits, which are symmetrical with respect to each atom (Figure 10.3).

The main thing is that the potential barrier is intact, the shape of which is no different from that of the internal pits. Schockley found the spectrum of electron energies for such a circuit and studied how this spectrum changes when the lattice constant shrinks from an infinite to a small size. The energy spectrum will be in the form of discrete (ring-ring) layers for individual atoms, and each line will be turned eight times. As the lattice constant decreases, it breaks down into individual layers, with layers intersecting at interatomic distances smaller than d_1 . In this case, two lines appear in the forbidden region of the crystal, one in the lower permissible region and the other in the upper region (Fig. 10.4).

Shock states occur as a result of the intersection of energy fields when the grid constant is small and the constant period is potential. Tamm states, on the other hand, arise as a result of changes in the shape of the periodic potential in the peripheral elementary layers.

Tamm and Shock-type surfaces can exist on a regular crystalline surface. On the actual surface of a semiconductor crystal obtained according to a particular technology, energetic states for electrons other than these states can also occur. For example, the ingress atoms deposited on the surface, the surface defects of the crystal lattice, dislocations and hakoses are caused by these. One of the important causes of surface states is adsorption (adhesion) of atoms foreign to the surface, in which the energy situation of the surface surface depends on the parameters of the lattice, as well as on the adsorption matter (adsorbant). Naturally, the density of surface surfaces in these cases will be different. When density is large, these layers can also form energy sectors.

10.2. Spatial charge layer on the surface of semiconductors

The presence of energetic layers on the surface of the crystal leads to the formation of an electrical charge in them. The true surface of the crystal can be recharged through the formation of an oxide layer in it, surrounded by medium, adsorbing ions.

The free charge carriers are attracted from the crystal volume to its surface area. Due to this, a double charged layer is formed.

Suppose that each atom on a surface gives a single surface surface. Let these layers be charged, in which case their density would be approximately 10^{15} cm^{-2}

Density of free electrons per unit of volume in metals 10^{22} cm^{-3} , in an atomic plane 10^{14} cm^{-2} will be in the coma. In order to reduce the surface charge on the metal, 10 atoms of the plane, i.e. 10^{-7} cm It is enough for the free charges in the step to

participate. This volumetric charge layer has no significant effect on the metal properties.

In semiconductors, the density of free charge carriers is much smaller than in metal. For example, if the density of free electrons in an n-type semiconductor is $n=10^{15} \text{ cm}^{-3}$ or one atomic plane is $N_1=10^7 \text{ cm}^{-2}$, then if the density of the surfaces is $N_s=10^{11} \text{ cm}^{-2}$, then to neutralize the surface charge is 10^{-4} cm . Free charge carriers will have to be removed from the layer in thickness, i.e., from the layer containing the atomic plane 10^4 . This layer is called the spatial charge layer. In semiconductors where the density of free charge carriers is even lower, the spatial (volumetric) charge layer can penetrate much infarther.

Thus, semiconductors have a surface spatial charge layer that can influence the properties of semiconductors. In this area, the energy sectors are bent. On the ordinate axis, the potential energy of the electrons in the crystal is the distance from the surface to the crystal on the axis of the absissa (Fig. 10.5).

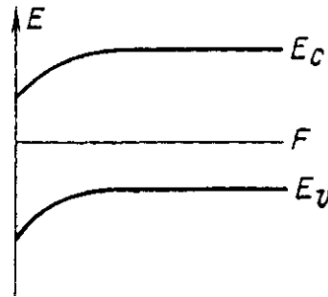


Fig. 10.5. Energy diagram for semiconductor surface

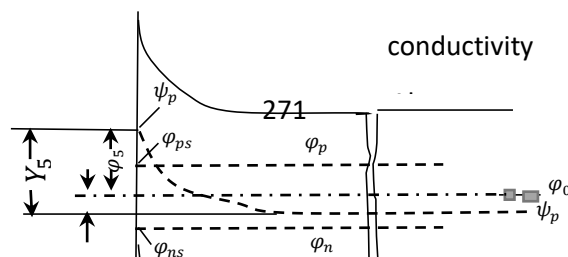
The depth of the spatial charge field is assumed to be equal to the Debay shielding length L for a private semiconductor :

$$L = \sqrt{\frac{\epsilon kT}{2\pi e n_i}} \quad (10.1)$$

In this case, the ϵ -dielectric constant, n_i is the density of charge carriers in a private semiconductor, and e is the electron charge.

A spatial charge field can form near the surface of a semiconductor. This field in the crystalline influences many properties of semiconductors. Therefore, the spatial charge layer of the semiconductor surface is studied.

The general theory of the spatial charge field was developed by Garrett and Brattain. Studies based on this theory made it possible to obtain quantitative information about the electron structure of the semiconductor crystal surface. Fig. 10.6 shows an energy diagram of a field near the n-type semiconductor surface.



$$\sim 10^{-4} \text{ SM}$$

Fig. 10.6. Energy diagram of a field near the surface of an n-type semiconductor

where, φ is the electrostatic potential, expressing the bending of energy fields: φ_0 is the potential in volume, which is assumed to be equal to the potential corresponding to the middle of the forbidden sphere; the surface value of this potential is $-\varphi_s$; F_o is the Fermian surface of the semiconductor volume; Y_s is the magnitude of the bend of energy fields on the surface. The distance from the electrostatic potential level on the surface to the Fermi level F_s is sometimes called the surface potential.

This theory poses the following questions:

- 1) determine how the spatial charge $\varphi(x)$ potential varies;
- 2) the density of excess charge carriers relative to volume in the near-surface area,
- 3) Find the magnitude of the volumetric charge and the surface $\Delta\delta_s$ conductivity they detect.

In theory, these hypotheses are adopted:

- 1) The density of inputs in the semiconductor volume is small, consequently, the semiconductor is undecomposed and the classical Bolsman statistics can be applied to electrons and pits;
- 2) the curvature of the spheres near the surface is small enough, in this place the semiconductor is also uneven, the boundary of the spheres near the surface is far enough from the Fermi level;
- 3) Inlets in crystal size are considered to be fully ionized.

In silicon, the level of the inlets of group III and V elements is fully ionized at room temperature (300 K). To find the law of change of electrostatic potential $\varphi(x)$ near the surface, it is necessary to solve the Poisson equation:

$$\frac{d^2\varphi}{dx^2} = -\frac{4\pi}{\varepsilon}\rho(x), \quad (10.2)$$

In this case, ε is the dielectric constant, $\rho(x)$ is the volumetric charge density.

Dimensionless magnitudes are entered:

$$Y = \frac{e}{kT}(\varphi - \varphi_0) = \beta(\varphi - \varphi_0), \lambda = \left(\frac{p_0}{n_0}\right)^{1/2}, P = \beta(F_V - F_0), N = \beta(F_n - F_0),$$

In this, $\beta = \frac{e}{kT}$, F_V , F_n are the Fermi quasi-levels.

The Poisson equation is much more complicated to solve. Equation (10.2) looks like this after some calculations:

$$\frac{dY}{dx} = \pm \frac{2}{L} \Phi(Y, \lambda, P, N). \quad (10.3)$$

In which $\Phi(Y, \lambda, P, N)$ is a particular function of the arguments in parentheses:

$$\Phi(\cdot) = \left[\lambda e^V (e^{-4} - 1) - \lambda^{-1} e^{-N} (1 - e^{+4}) + (\lambda - \lambda^{-1}) Y \right]^{1/2}. \quad (10.4)$$

The bending of spheres near the surface from the energy scheme of the surface:

$$Y_s = \beta(\varphi - \varphi_0) = \beta F_s + \ln \lambda, \quad (10.5)$$

It can be expressed in appearance. For a private semiconductor, $\lambda = 1$ and $Y_s = \beta F_s$. The more a semiconductor differs from specificity, the more $Y_s = \beta F_s$ differs from each other. The conductivity on the surface of the semiconductor, the density of the excess charge carriers, and other magnitudes are considered in dependence on βF_s .

(10.3) tenglamadan

$$x = \frac{1}{2} L \int_0^{Y_s} \frac{1}{F(Y, \lambda, P, N)} dY, \quad (10.6)$$

origin. If the integration in the equation were performed, the dependence of the surface potential φ on x would be found. But this calculation is much more complicated. Instead, it is more convenient to find a function of $\varphi(x)$ in several separate areas of the spatial charge layer. Three spheres are usually considered: the thinned sphere, the inversion layer, and the enriched sphere.

1. The density of free charge carriers in a thinned field will be much smaller than the density of uncompensated inputs. If in an n-type semiconductor the spheres are bent upwards ($Y_s < 0$), then in a p-type semiconductor the spheres are bent downwards ($Y_s < 0$), then a thinned field is formed.

In this area $F(Y, \lambda, P, N)$ in the function $F(Y, \lambda, P, N)$ the first and second hads are smaller than the third, consequently:

$$\frac{dY}{dx} = \pm \frac{2}{L} (Y, \lambda - \lambda^{-1}) \frac{1}{2} Y \frac{1}{2}. \quad (10.7)$$

The positive sign is for cases where the spheres are bent upwards and negative ones are for cases where the spheres are bent down. (10.7) After the expression is integrated, the following expression

$$Y = Ax^2 \quad (10.8)$$

origin. However, this is a narrow area in which the law is enforced (Fig. 10.7, a, b).

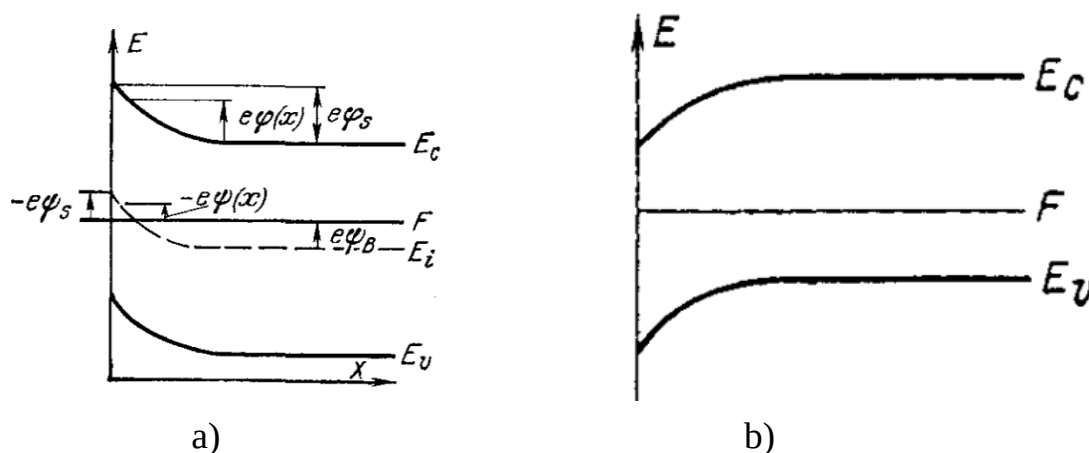


Fig. 10.7. Energy diagram of the semiconductor surface domain

2. Now let's look at n-type semiconductors. In the case of a large upward bend of the spheres, let the sphere approach the Fermi level relative to the conductivity area.

In this layer, the conductivity of the cavity may be greater than the electronic conductivity. Such a layer is called an inversion layer. To find the patterns of potential variation in this layer $F(Y, \lambda, P, N)$ Taking into account only one addition, either the first or the second addition, equation (10.3) is solved. This choice depends on the type of conductivity of the semiconductor. The bending value of the spheres depends on the degree of lequefaction of the sample (the density of electron-giving or hollow-forming inlets). Generally speaking, the equation (10.3) for the inversion layer is solved in the same way.

3. If the spheres in the n-type semiconductor are bent downwards, and the spheres are bent upwards in the p-type semiconductor, then an enriched field is formed, i.e., the density of the main charge carriers is greater near the surface relative to the volume, but the width of the enriched sphere is also small enough, as in the case of an inversion layer. By throwing out the third case in the function $F(Y, \lambda, P, N)$ and solving the equation (10.3), the potential distribution $\phi(x)$ is found. In general, by calculating Equation (10.3), it is possible to obtain a bond of $\phi(x)$ over the entire spatial charge region.

10.3. Field effect in semiconductors

A more perfect way to change the potential on a surface is through a

phenomenon called the field effect. It consists in the following: if one coating is a sample and the other is metal, a large dielectric absorption between the coatings (ϵ) If a capacitor on which an insulator (dielectric) is located is assembled and voltage is applied to it, in this case a volumetric charge is generated near the surface of the semiconductor. It neutralizes the charge on the metal plate.

Suppose that the initial bending of the energy fields for this (p-type) conductor was zero until the field (voltage) was inserted (Fig. 10.8-a). When voltage is given (Fig. 10.8b), the density of the canopies in the area near the surface is greater than the volume, while the energy fields bend upwards.

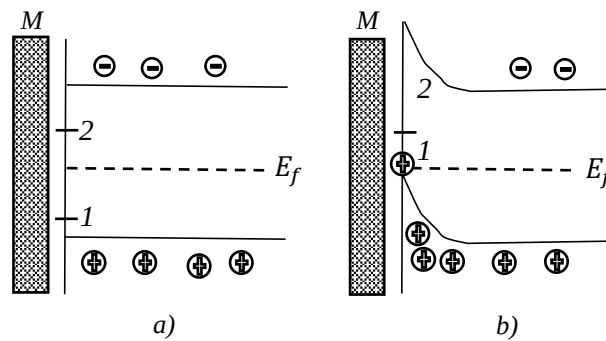


Fig. 10.8. (Initial a) and (bent b) states of energy fields for a p-type conductor

It is possible to modify the potential potential at the surface of the semiconductor sample by varying the area to be generated at the surface. In this change, the distribution of charges in the surface area changes. The phenomenon of change in the crystalline conductivity of a semiconductor under the influence of an electric field is called the field effect.

As we know, there will be electron states on the surface of the semiconductor. Some of the charges generated by the action of the field sit in these states. The greater the density of cases at the surface, the smaller the observed stationary change of conductivity in the field effect phenomenon. Therefore, field effect measurement can be used to gain information about the parameters of states on the surface. When field effect tests are carried out, the system of states on the surface does not change. In this case, only the saturation of the cases on the surface changes. However, when too much voltage is created, the phenomenon of electroadsorblation occurs. With the field effect method, it is possible to change the potential on the surface practically continuously at a sufficient interval and generate the desired areas of curvature.

Based on the measurement of the field effect in the stationary (time constant) state, the type of conductivity on the surface (*n-tur* or *P-trip* It is possible to find the

initial bend of the spheres on the surface and determine whether the bending of the spheres depends on the applied voltage. Based on the field effect data, it is possible to determine the amount of charge in the surface area and the amount of charge that passes to the surface depending on the bending of the spheres. Finally, the study of field effect kinetics can help determine the reaction time of states on the surface.

This method gives good results in cases where the electron state density on the surface is not very large (10^{10} - 10^{12} cm^{-2}) Let's assume that there is an equilibrium between areas of volume and cases on the surface. In the experiment, the bonding of the sample conductivity $\Delta\sigma$ change to the voltage $V\Delta$ put into the capacitor $\Delta\sigma=f(V)$ or the total charge generated by the field to Q is measured $\Delta\sigma = \phi(Q)$ i.e., $Q = CV$, wherein C is the capacitor capacity, usually having a value of a few decimal times C .

Figure 10.9 (n-Ge), Fig. 10.10 (n-Si) represents the curves of the field effect obtained in the experiment. Below are the formulas for volumetric charge Q_s , surface field maydon E_s and electrons, and pit densities Γ_n and Γ_p :

$$Q_s = 2en_iLF(Y_s, U_b), \quad (10.9)$$

$$E_s = \frac{2en_iL}{\varepsilon} F(Y_s, U_b), \quad (10.10)$$

$$\Gamma_p = n_iLG(Y_s, U_b), \quad \Gamma_n = n_iLG(-Y_s, U_b), \quad (10.11)$$

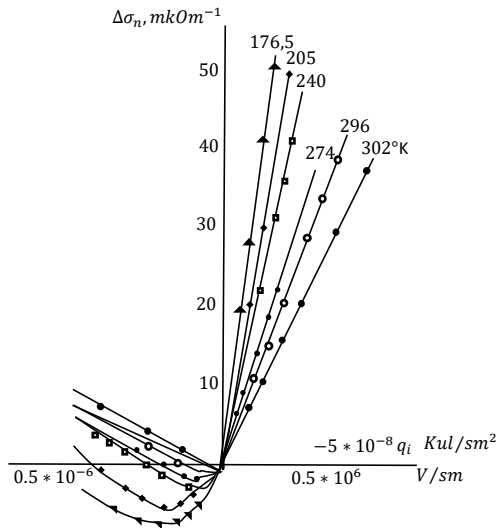


Fig. 10.9. Of the field effect for *n-Ge*
Curves obtained in the experiment

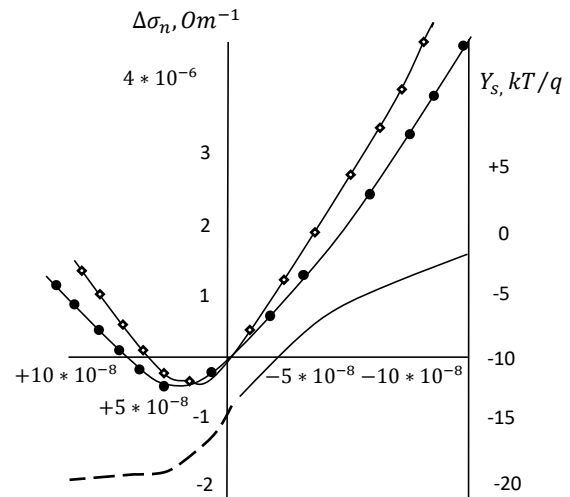


Fig. 10.10. Experimental Curves of Field Effect
for *n-Si*

wherein the gesture of change of conductivity on the surface determines a value of this type of conductivity. The binding minimum of surface conductivity to the

bending of spheres depends on the volumetric properties of the crystal and does not depend on the initial bending of the spheres. Therefore, the minimum is the zero point, from which are the experimental and theoretical values of conductivity. When calculating the theoretical curve for a material with a full ionized input in volume, it is necessary to know the density and mobility of the charge carriers in the volume.

The volumetric charge at the depth of shielding Q_s , the area on the surface is E_s and the electrons and the densities of the spheres are Γ_n and Γ_p , and if the spheres are bent Y_s , in these expressions $U_b = \ln \lambda / \beta$ - is the distance in kT/e units between the Fermi surface and the sphere prohibited in volume. Permeability at the surface

$$\Delta\sigma = \varepsilon(\mu_p \Gamma_p + \mu_n \Gamma_n), \quad (10.12)$$

is determined by expression. The calculation becomes much more complicated when the degree of ionization of inputs in the volume is optional.

10.3.1. Determination of surface parameters based on field effect measurements

The external electric field (the inlet of the field) dissipates the charge Q_{sc} at the depth of the shielding. Total charge generated by the action of the field

$$Q = Q_{sc} + Q_{ss} \quad (10.13)$$

For each voltage value placed on the electrode (metal), the magnitude of Q can be calculated on the basis of $Q = CV$. From the connection $Y_s(V)$ it is possible to determine the bending of the fields corresponding to this field. For found is Q_{sc} . Charge on the surface

$$Q_{ss}(Y_s) = CV - Q_{sc} Y_s \quad (10.14)$$

is found in the expression.

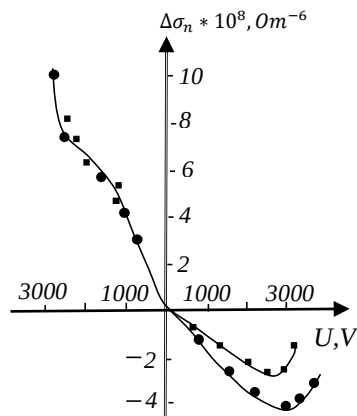


Fig. 10.11. (σU) bond obtained for German

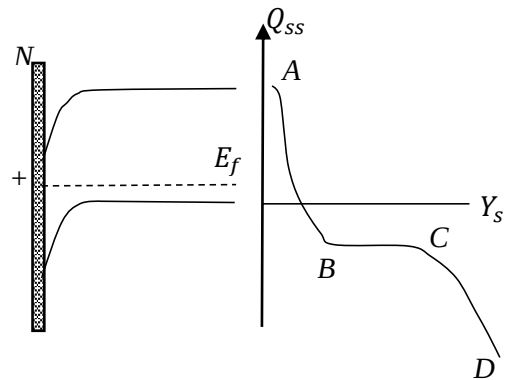


Fig 10.12. Received link for German $Q_{ss}(Y_s)$

Also, the link obtained for the Germanic is represented in Fig. 10.11. In a number of cases, $Q_{ss}(Y_s)$ bonding determines the discreteness or continuity of the layers on the sample surface. If the surfaces are discrete and located at a sufficiently large distance from each other, the $Q_{ss}(Y_s)$ curve should be rung. This can be seen in Fig. 10.12. When the sphere is said to bend downwards ($Y_s > 0$) the l-surface near the Fermi level is initially filled (Fig. 10.12- b, Part AB).

At a Y_s value, it is fully filled, but the 2-nd line is far from the Fermi level. As Y_s increases (as the sphere bend increases), the Q_{ss} charge remains constant in the BC part, and then (*the CD* part) changes to the 2-nd line filling. Such a state of affairs has been observed for CH_2O .

For silicon and germanium, $Q_{ss}(Y_s)$ the changing of the world is without a slippery slope. Such a change:

- 1) when the energy levels on the surface have a continuous spectrum;
- 2) At surface surfaces located close to each other, this position becomes appropriate.

At low temperatures, discrete layers are observed in Germanium. Field effect measurement results to a mobility-dimensionable magnitude

$$M_{me} = \frac{d\Delta\sigma_s}{dQ} \quad (10.15)$$

It is convenient to describe using it. This is called field effect mobility. For it, the following theoretical expression was obtained:

$$\mu_{me} = \mu \frac{b^{-1} - \lambda^{-2} \exp Y_s}{1 + \lambda^{-2} \frac{Z}{\exp Y_s} \varphi \exp \varphi_s}, \quad (10.16)$$

In this

$$Z = \frac{2F}{\lambda n_i L}; \quad \varphi = \frac{\exp \varphi_s}{e} \frac{dQ_{ss}}{d\varphi_s}. \quad (10.17)$$

For silicon μ_{ep} , the parameters of surface surfaces φ_p were measured and the curve of adhesion to surface surfaces was analyzed in a simple way.

10.3.2. Field effect kinetics

In early studies of the field effect, it was confirmed that surface conductivity changes rapidly when a field is formed, and then first very quickly, then slowly, decreasing to a certain stationary value (sometimes to zero). Such a change in the field effect has been observed in germanium, silicon, Cu_2O , PbS , CdS , $GaAs$ semiconductors. Fig. 10.13 shows Δ the binding of $\Delta\sigma_s$ to τ . As can be seen from the figure, at the time of the order of 10^{-8} s, the conductivity on the surface will be

much lower than its original value, while at 10^{-5} s it will have some *kind of* $\Delta\sigma_s$ quasi-relative value.

Then another 10^{-3} s seconds to 1 second, even up to hours, the stationary value takes over the $\Delta\sigma_{sbar}$.

The reason why $\Delta\sigma_s$ decreases over time is that some of the charge carriers that form a field pass to the surfaces. He can no longer participate in the playoffs. This property of field effect kinetics confirms that there are two systems of layers in germanium and silicon, their relaxation time differing strongly from each other. $\tau_s \leq 10^{-8}$ s Those that have a relaxation time are called "fast" layers, and $\tau_s \geq 10^{-3}$ s those that have a relaxation time are called "slow" layers.

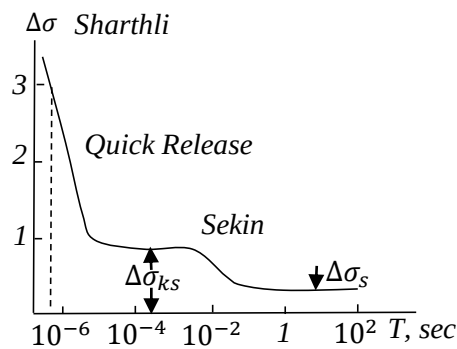


Fig. 10.13. Δ Binding of σ_s to τ

This large difference is related to the location of layers. It is known that on the actual surface of germanium there will be an oxide layer of $\sim 10 \text{ \AA}$ thickness. In air, the thickness of this layer increases to $\sim 50 \text{ \AA}$, after heat treatment it can become even thicker. Cases on the outer surface of the oxide are "slow" states because charge carrier exchange with spheres of their volume is difficult, and the relaxation time is large.

The states on the surface located at the oxide-semiconductor boundary become "fast" states because their τ_s are small. We take a closer look at the field effect kinetics using the diagrams in Fig. 10.14.

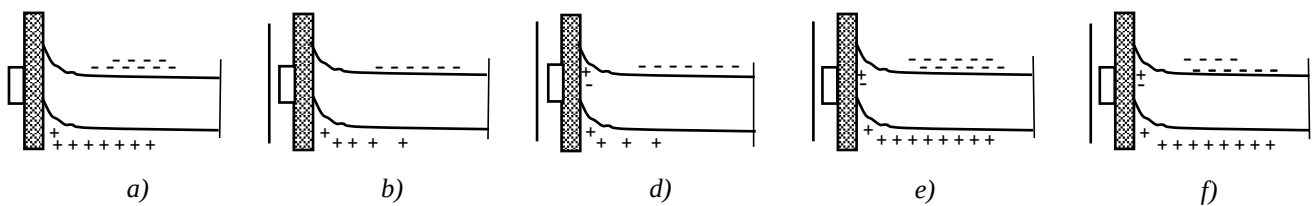


Fig. 10.14. Field effect kinetics diagrams

Let an inversion layer be present on the n-Ge sample surface until it is inserted into the field (the initial bend of the spheres is facing upwards, Fig. 10.14).

We insert an area so that the metal plate becomes negative (we give voltage) (Fig. 10.14-b). To neutralize the charge on the plate, non-basic charge carriers (kovakiars) go from the volume to the area near the surface, and the electron to the volume metal electrode.

The height of the potential barrier near the surface varies in such a way that in the volumetric charge region there are N surplus pits, and in the volume there are N fewer electrons and pits. Consequently, the sample conductivity varies:

$$\Delta\sigma_s = eN(\mu_{ps} - \mu_{nb} - \mu_{pb}). \quad (10.18)$$

It occurs during the resettlement process of $\Delta\sigma_s$ during the formation of a diffusion-drift equilibrium (Maxwell relaxation time) $\tau_m = \frac{\varepsilon}{4\pi\sigma_0}$ in which σ_0 is the

specific resistance of the sample. This time is true to the time of diffusion of the charge carriers to a distance equal to the shielding distance and is determined by the RC of the material in question. For the inspected Ge samples, it was $\sim 10^{-11} \sim 10^{-12}$ s

After that, an equilibrium arises between the spheres and the "fast" layers during the relaxation time $\tau_s - 10^{-8}$ s (Fig.10.14-d): the total charge in the "fast" surface and the volumetric charge area is equal to eN . If we define the portion of the total charge located on the surface by x , then there is a change in conductivity 10^{-8} s after the field is generated

$$\Delta\sigma_s = eN[(1-x)\mu_{ps} - \mu_{nb} - \mu_{pb}] \quad (10.19)$$

bo'ladi.

In the sample, when the overcharge carriers have a residence time (of $\tau \sim 10^{-5}$ s for Ge), an equilibrium of the densities of the swells and electrons that exist until the field is introduced is established (Fig. 4.7e).

Variation of permeability after $\sim 10^{-5}$ s

$$\Delta\sigma_s = eN(1-x)\mu_{ps} \quad (10.20)$$

$\Delta\sigma_s$ The next change in is due to the transition of charge carriers to "slow" surface layers (Fig. 10.14- f). Since their density is large, all charges generated by the field can pass to "slow" layers. The bending of the spheres near the surface returns to its previous value, while $\Delta\sigma_s$ tends to zero. This process occurs during the "slow" state relaxation, which varies in a wide interval (from minutes to hours).

As we know, a number of important parameters are obtained based on the investigation of the kinetics of the field effect at different time intervals. "Fast" and "slow" surface levels are found to be periods of relaxation and residence time for excess charge carriers.

In this case, the surface mobility of the charge carriers, which is equal to the field effect mobility, can also be detected.

10.4. Surface recombination rate

In most semiconductors, recombination of charge carriers occurs mainly through input centers (recombination centers). They will be located near the middle of the prohibited sphere in volume. Such a theory of recombination was developed by Shockley and Reed. There is also the possibility of inter-zone recombination, but its probability is much smaller.

In pure crystal sizes there are few recombination centers so recombination via surface surfaces may be important. In thin samples, electrons and swells are mainly recombined on the surface, and therefore the effective residence time of charge carriers is determined by this recombination. Surface electronic states perform the functions of recombinatory centers on the surface. They bond well with spheres of volume, and have sufficiently large cross-sections to trap charge carriers.

The study of surface recombination processes provides information about the nature of surface layers, their energy system. Regardless of some specifics, surface and volumetric recombination are basically subject to the same laws. However, in some cases it is more convenient to study surface recombination than to know volumetric recombination. Such a correlation of the surface recombination rate to the electron uptake of the recombinant state (i.e., the state of that state relative to the Fermi level) can be obtained in a relatively simple way.

The theory of surface recombination was developed on the basis of the Shockly-Reed theory.

Let there be a kind of recombination centers on the surface with an energy level N_s density and E_s . The excess charge carriers generated in volume and having Δn density are recombined by the centers on the surface. We calculate the same recombination rate.

If we define the probability of an electron in the state E_s by f_s , then its expression of the thermodynamic equilibrium state would be:

$$f_{0s} = \left(\exp \frac{-E_s - E_F}{kT} + 1 \right). \quad (10.21)$$

The rate of transition of electrons from the conductivity region to the E_s state is proportional to the density n_s near the surface, and E_s to the unoccupied part of the surface:

$$-\left(\frac{dn_s}{dt} \right)_1 = \chi_{ns} n_s N_s (1 - f_s), \quad (10.22)$$

In this case, χ_{ns} is the coefficient of the electron to be in the E_s state, the

velocity of departure of electrons from this state:

$$\left(\frac{dn_s}{dt}\right)_2 = \beta_{ns} f_s N_s, \quad (10.23)$$

Wherein β_s is the current transition coefficient. Clearly, the electron density change is

$$-\left(\frac{dn_s}{dt}\right)_1 = -\left(\frac{dn_s}{dt}\right)_1 - \left(\frac{dn_s}{dt}\right)_2 = \chi_{ns}(1-f_s)n_s N_s - \beta_{ns} f_s N_s, \quad (10.24)$$

Under equilibrium:

$$\chi_{ns}(1-f_s)n_s N_s = \beta_{ns} f_s N_s. \quad (10.25)$$

Using the expression (10.21),

$$\beta_{ns} = \chi_{ns} N_s \left(\frac{-E_s}{kT}\right), \quad (10.26)$$

Now

$$\begin{aligned} \beta_{ns} &= \chi_s \exp\left(\frac{-E_s}{kT}\right) \\ -\frac{dn_s}{dt} &= \chi_{ns} n_s N_s [(1-f_s)n_s - n_{1s} f_s], \end{aligned} \quad (10.27)$$

In this

$$n_{1s} = N_s \exp\left(\frac{-E_s}{kT}\right). \quad (10.28)$$

In the same way we find the huge variations in the density of the canines:

$$-\frac{dp_s}{dt} = \chi_{ns} N_s [p_s f_s - p_{1s}(1-f_s)], \quad (10.29)$$

In this

$$p_{1s} = N_s \exp\left(\frac{E_s}{kT}\right). \quad (10.30)$$

The electrons and canine density changes are equal to each other as a result of recombination, i.e., $\frac{dn_s}{dt} = \frac{dp_s}{dt}$ and from this fs are found:

$$f_s = \frac{\alpha_{ns} n_{1s} + \alpha_{ps} p_{1s}}{\alpha_{ns}(n_s - n_{1s}) + \alpha_{ps}(p_s + p_{1s})}. \quad (10.31)$$

The electron density change as a result of recombination is now

$$-\frac{dn_s}{dt} = \frac{\alpha_{ns} \alpha_{ps} N_s (n_s p_s - n_{s0} p_{s0})}{\alpha_{ns}(n_s - n_{1s}) + \alpha_{ps}(p_s + p_{1s})} \quad (10.32)$$

$\frac{dn_s}{dt} = \frac{\Delta n_s}{\tau_s} = s$ Surface refers to the rate of recombining, that is, $1/s$ of the surface in 1 cm^2 refers to the number of electrons and holes that are recombined on the surface, and cm/s has a unit. Volumetric R vs n Surface with densities p_s and n_s Links Between Densities:

$$p_s = p e^{\frac{\phi_s}{kT}}, n_s = n e^{-\frac{\phi_s}{kT}}, \quad (10.33)$$

As in appearances, $n_s \cdot p_s - n_{s0} \cdot p_{s0}$ separation $p \cdot n - n_0 \cdot p_0$ can be replaced by:

$$p = p_0 + \Delta p, n = n_0 + \Delta n, \Delta n = \Delta p_s. \quad (10.34)$$

The result is the following expression for s :

$$s = \frac{(p_0 + n_0 + \Delta n) Y_{ns} Y_{ps} N_s}{Y_{ns} n_1 + Y_{ps} p_1 + [Y_{ns} n \exp(-\phi_s / kT) + Y_{ps} p \exp(\phi_s / kT)]} \quad (10.35)$$

In the case of weak generation Δ , $p \ll p_0$, $\Delta n \ll n_0$, but $\Delta p = \Delta n$, and the rate of recombination is

$$s = \frac{(p_0 + n_0) Y_{ns} Y_{ps} N_s}{Y_{ns} (n_s + n_1) + Y_{ps} (p_s + p_1)} \quad (10.36)$$

It looks simple. here

$$p_1 n_1 = n_i^2, p_s n_s = pn. \quad (10.37)$$

Is there a limit to the application of the concept of surface recombination rate? Oh, yes. We will see it in the analysis below. A quantitative study of this issue leads to the following conclusions:

1. If the generation across the whole volume is uniform and the surface potential is not very large, the concept of surface recombination rate introduced above is fair.
2. If the excess charge carrier generation occurs only in the surface region of the shielding thickness L , then $p_s n_s = pn$ Equality is not enforced, that is, the equilibrium between volume and surface is broken.
3. length of diffusion of a charge carrier l_d spatial charge domain L_s If it is smaller than the length of the shell, in this case the recombination of the same layer becomes significant, the volume and surface balance is disturbed.

According to the results of the calculations, $p_s n_s = pn$ the condition that equality is fulfilled

$$L_1 = \left(L_2 + \frac{s}{v_d} \right); \quad (10.38)$$

$$L_1 = \frac{1}{l_d} \int_0^p e^{-y} dx; L_2 = \frac{1}{l_d} \int_0^1 e^y dx; v_d = \frac{l_d}{\tau_p}, \quad (10.39)$$

In this, l - the distance from the surface to the plane where $Y = 1$; l_d is the length of diffusion shear; τ_p - Residence time of hives in electronic semiconductors; v_d is the diffusion shear rate.

(10.38) has been confirmed to satisfy the condition in the field of bending of the spheres used in practice for germanium. But big enough Y_s . And it's huge s . When recombination rates lose their meaning. High Om li crystals (past Om (10.38) may not be performed at small bends (10.38). The concept of s should be used with caution in surface potency values, which we use in practice for silicon. Again, the concept of surface recombination velocity in semiconductors, which is as wide as the forbidden field, cannot be applied to the bending of very large sub-spheres.

10.5. Effect of surface recombination on the residence time of poison carriers in finite size samples

The presence of surface recombination traps naturally leads naturally to increasing of the total number of recombination sites in the sample and consequently to a shorter residence time for charge carriers.

Consider the distribution of the excess concentration of charge carriers near the surface, where $s \neq 0$ the formation of excess electron-pit pairs occurs uniformly over the entire sample volume at the speed of G .

For precision we select an n-type semiconductor whose x-coordinate is normal to the sample surface and lifts as it penetrates from the surface into the depth of the sample, and there is no external electric field.

Non-primary carriers - excessive concentration of hives $p\Delta(x)$ can be found in solving the stationary continuity equation

$$\frac{\partial p}{\partial t} = -\frac{1}{e} \frac{\partial J_p}{\partial x} - \frac{\Delta p}{\tau_p} + G = 0. \quad (10.40)$$

Since the strength of the electric field in the field of spatial charge is small, the drift organizer of the current cannot be taken into account. Without it

$$J_p = -eD_p \frac{dp}{dx}$$

and (10.40) the equation takes the form as follows:

$$D_p \frac{d^2 \Delta p}{dx^2} - \frac{\Delta p}{\tau_p} + G = 0. \quad (10.41)$$

The boundary condition for the flow of hives to the surface at the point $x = 0$ can be given as follows:

$$D_p \frac{d\Delta p}{dx} = -\frac{J_p}{e} = \Delta q_s = s\Delta p. \quad (10.42)$$

Since the electron and pit fluxes resulting from surface recombination are of pure diffusion nature, it can be expected that the effect of surface on the distribution of excess pit concentration will be expected to be felt at distances on the order of several diffusion lengths from the surface. So the continuity equation for $x \rightarrow \infty$ (10.41) is:

$$-\frac{\Delta p}{\tau_p} + G = 0 \text{ and } \Delta p|_{x \gg L_p} = G\tau_p. \quad (10.43)$$

The general solution of (10.41) looks like this:

$$\Delta p(x) = Ae^{-x/L_p} + G\tau_p. \quad (10.44)$$

To find A, we use the boundary condition (10.42) at the point $x = 0$:

$$-\frac{A}{L_p} D_p = s(A + G\tau_p). \quad (10.45)$$

here

$$A = -\frac{Gs\tau_p L_p}{D_p + sL_p} = -\frac{Gs\tau_p^2}{L_p + s\tau_p}. \quad (10.46)$$

Finally, we get the following statement:

$$\Delta p(x) = G\tau_p \left[1 - \frac{s\tau_p}{L_p + s\tau_p} e^{-x/L_p} \right] = G\tau_p \left[1 - \frac{s/v_d}{L_p + s/v_d} e^{-x/L_p} \right], \quad (10.47)$$

This $v_d = L_p / \tau_p$ is where the diffusion rate.

The physical meaning of the resulting relationship is very clear:

1) If $s \ll v_d$ there is, then small gradients of concentration of excess charge carriers are needed for surface recombination to provide a small diffusion flux, i.e. with the change x Δp changes slightly, and $\Delta p \approx G\tau_p$;

2) $s \gg v_d$ At dp is strongly dependent on x , since there must be sufficient diffusion flux for surface recombination. The nature of the dependence on ni x in different ratios of s and v_D $\Delta n = \Delta p$ is illustrated in Fig. 10.15.

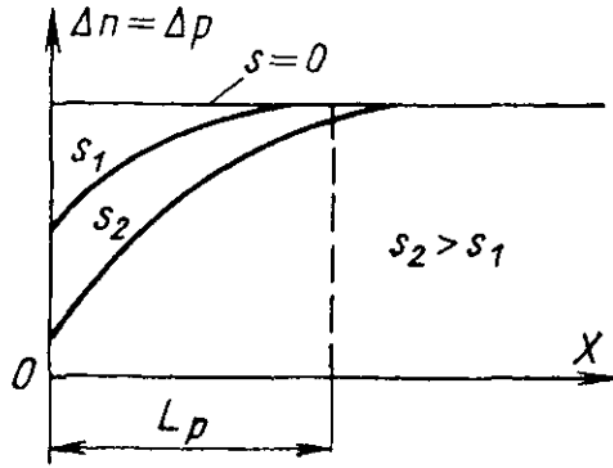


Fig. 10.15. Effect of surface recombination rate on the distribution of concentration of excess charge carriers

It is possible to determine the concentration of Δp plus p-fins on the semiconductor surface by (10.47) by putting $x=0$:
equals:

$$\Delta p_s = \frac{G\tau_p L_p}{L_p + s\tau_p}. \quad (10.48)$$

$\tau_s = \Delta p_s / G$ This is the residence time of non-primary charge carriers on the surface of the semiconductor. From (10.48) comes the following expression:

$$\tau_s = \tau_p \frac{L_p}{L_p + s\tau_p}. \quad (10.49)$$

Without it

$$1 - \frac{\tau_s}{\tau_p} = \frac{s\tau_p}{L_p + s\tau_p}. \quad (10.50)$$

In this case (10.47) the parity is as follows:

$$\Delta p(x) = G\tau_p \left[1 - \left(1 - \frac{\tau_s}{\tau_p}\right) e^{-x/L_p} \right] = G\tau_p \left[1 + \left(\frac{\tau_s}{\tau_p} - 1\right) e^{-x/L_p} \right]. \quad (10.51)$$

From here we get the residence time of non-primary charge carriers in the near-surface region of the semiconductor:

$$\tau(x) = \frac{\Delta p(x)}{G} = \tau_p \left[1 + \left(\frac{\tau_s}{\tau_p} - 1\right) e^{-x/L_p} \right]. \quad (10.52)$$

Control Ceremonies

1. Tell us about the nature of surface surfaces
2. Explain the Kronig-Penny model of weakly bonded electrons

3. Describe the Tamm and Shock types
4. What do you mean by the spatial charge layer on the surface of semiconductors?
5. Explain the field effect in semiconductors
6. What is it like to determine the parameters of surfaces based on field effect measurements?
7. What do you mean by field effect kinetics?
8. Explain the rate of surface recombination
9. Explain the Shockly-Reed theory
10. What will be the effect of surface recombination on the residence time of venomous carriers in finite size samples?

Chapter XI

ABSORPTION OF LIGHT IN SEMICONDUCTORS

11.1. Return spectrum and absorption spectrum

Suppose that a monochromatic radiation with intensity I_0 falls on a semiconductor plate (Fig. 11.1). Due to the process of reflection and absorption, the intensity of light passing through the sample decreases. If I_R is the intensity of the reflected light, then the proportion of reflected radiation is assigned to the following ratio

$$R = I_R / I_0 \quad (11.1)$$

that is, it will be equal to what is called the reversal coefficient.

The dependence of the reflection coefficient of matter on the energy of the quantum of falling light $R(h\nu)$ or $R(\lambda)$ on the wavelength is called the reversal spectrum.

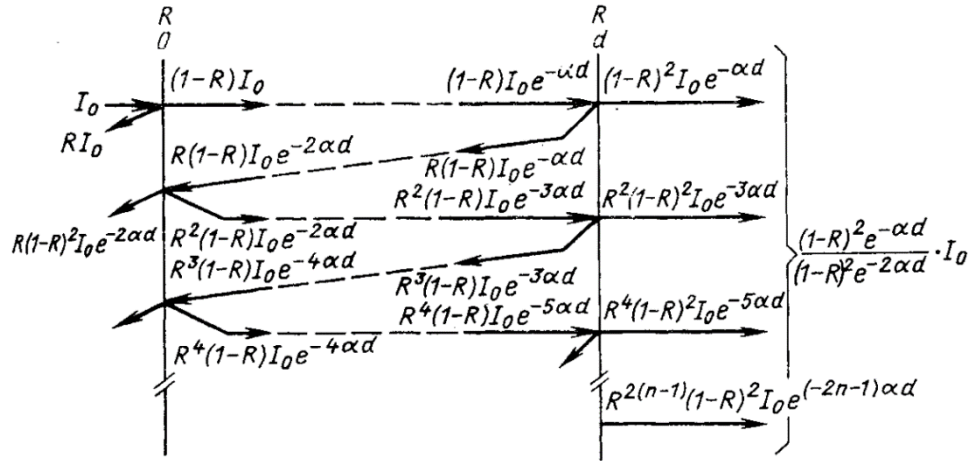


Fig. 11.1. Flow of energy in a multi-internal feedback system

Radiation considering bounce $(1-R) I$ Pass through the first surface of the specimen. dx In the unit of time due to the absorption of light in the layer I Intensity of radiation dI decreases. Amount of energy absorbed dI The amount of energy going into the layer I and the thickness of the adsorption layer DX ga proportionslindir:

$$-dI = \alpha I dx. \quad (11.2)$$

The coefficient of proportionality, which represents the amount of energy absorbed by a substance from the unit radiation intensity in a unit-length layer, is called the α absorption coefficient.

By integrating Equation (11.2), we find the intensity of the radiation that reaches the second surface of the plate $(1-R)I_0 e^{-\alpha d}$. The intensity of the light emitted from the sample is as follows $(1-R) \cdot (1-R)I_0 e^{-\alpha d}$: Fig. 11.1 shows that the light returning from the sample comes out of it attenuated. Taking these effects into account, the conversion coefficient T , d will be equal to the ratio of the intensity of

the light passing $I = \frac{(1-R)^2 e^{-\alpha d}}{1-R^2 e^{-2\alpha d}} I_0$ through the thickness sample to the incidence of the luminous intensity I_0 .

$$T = \frac{(1-R)^2 e^{-\alpha d}}{1-R^2 e^{-2\alpha d}}. \quad (11.3)$$

If αd is large, the second case in the sentence of (11.3) can be discarded. In this case, the intensity of the light passing through the d -thickness sample, taking into account regression, will be equal to:

$$I = (1-R)^2 I_0 e^{-\alpha d}. \quad (11.4)$$

Here α is characteristic of the absorbent medium that depends on the wavelength of the radiation. The dependence of the absorption coefficient on the energy of the quantum of light $a(h\nu)$ falling into matter or the wavelength $\alpha(\lambda)$ is called the absorption spectrum.

Let there be n absorption centers. We define the probability of one photon being absorbed by a single absorption center in a unit of time α . In other words, σ is the efficient absorption cross-section of a single photon per unit of time. Effective absorption depends on the transverse cross-section σ photon energy and the nature of the absorption centers.

If we define the mean free travel path of a photon in an absorbing medium by l_f , according to (5.7) it becomes equal to:

$$l_f = (\sigma N)^{-1}, \quad (11.5)$$

and absorption coefficient

$$\alpha = 1/l_f = \sigma N \quad (11.6)$$

This refers to the probability of absorption in the length of the photon.

Let's assume that a semiconductor has absorption centers of different natures. If the absorption centers N_i are characterized by an effective absorption section σ_i , then

$$\alpha_i(h\nu) = 1/l_{if} = \sigma_i N_i. \quad (11.7)$$

The coefficient of complete absorption of the substance α is as follows:

$$\alpha(h\nu) = \sum_i \sigma_i N_i = \sum_i \alpha_i, \quad (11.8)$$

That is, the absorption coefficients of the different centers are equal to the sum of them. Consequently, the complete absorption spectrum of a substance consists of the sum of the absorption spectra of the different absorption centers.

When the electrons of a semiconductor interact with electromagnetic radiation, two laws must be fulfilled: the law of conservation of energy and the law of conservation of quasi-impulse. If $h\nu = \hbar\omega$ before interacting with a quantum of light, which $\hbar\eta$ has an electron energy and an impulse, its energy is equal to E and

its quasi-impulse is equal to p , and has E' and p' respectively from the interaction, then these laws are written as follows:

$$E' = E + \hbar\omega; \quad (11.9)$$

$$p' = p + \hbar\eta. \quad (11.10)$$

Absorption of radiation in semiconductors can be conversed with a change in the energy state of the free or atomically coupled electrons, as well as with a change in the vibrational energy of the lattice atoms. In this regard, optical absorption in semiconductors is divided into five main types: private, exciton, absorption by free charge carriers, ingress absorption, and lattice absorption.

11.2. Private Defeat in Right Transitions

If a semiconductor absorbs a quantum of light, electrons are excited from the stress band to the superconductivity zone, then such a loss is called private or fundamental. When studying the private absorption of a semiconductor, it is necessary to take into account the structure of its energy zones. Semiconductors as we currently know can be divided into two main types, according to the energy zone configuration. In the first of these, the minimum energy in the conductivity zone, characterized by the wave vector $k_{\min}$, and the maximum energy in the stress zone, defined by the $k_{\max}$ wave vector, are located at exactly one point in the Brillien zone (usually at the point $k = 0$). In other words, it is $k_{\min} = k_{\max}$ for semiconductors (Fig. 11.2-a). An example of a semiconductor with such a zone structure would be indium antimonide (see Fig. 2.25).

For substances of the second type, the extremes of the conductivity and tensile zones are at different k , so in this case the $k_{\min} \neq k_{\max}$ is (Fig.11.2-b). The latter type of substances include most semiconductors, including germanium and silicon (see Fig. 2.21).

The passage of electrons through the bandgap occurs primarily between the energy states corresponding to the maximum of the stress band and the minimum of the conduction zone, i.e., at values close to zero of quasi-impulse p or wave vector k , as shown in Fig. 11.2-a.

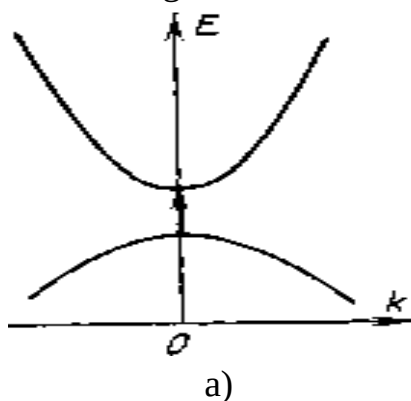
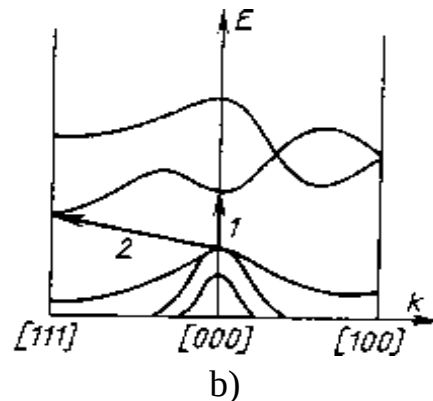


Fig. 11.2. Private absorption in interband straight transitions (a) and optical transitions in a complex energy-band semiconductor (b)

The photon impulse $\hbar\eta = 1/\lambda_f$ is very small in comparison with the electron quasi-impulse, since the wavelength of light is 10^{-1} - 10^{-5} cm and 300 K and the wavelength of an electron in a crystal corresponding to an electron energy of about is about $5 \cdot 10^{-7}$ sm. In this regard, the photon impulse can be ignored in (11.10), so

$$p^{\parallel} = p \quad (11.11)$$

or

$$k^{\parallel} = k. \quad (11.12)$$

The relation called the (11.11) or (11.12) electron transition selection rule indicates that there are such transitions in which the wave vector of the electron is conserved during the interaction of the crystal electron with the radiation field. These jumps are called vertical or straight jumps (Fig. 11.2-a and 1-jump in Fig. 11.2-b).

For a semiconductor whose conductivity band and stress band have spherical symmetry, consider private absorption during direct interband transitions. Suppose the tensile zone is completely occupied by electrons, and the conduction band is completely empty. This condition is appropriate for unalloyed semiconductors with a density of 0 K⁰. According to Fig. 11.3, each initial state E in the stress zone passes into the final E' state in the conductivity zone as a result of the absorption of a quantum of light $h\nu$ with energy:

$$E^{\parallel} = h\nu + E. \quad (11.13)$$

In the case of spherical symmetry zones according to (4.4) and (4.10):

$$E^{\parallel} = E_C + \frac{\hbar^2 k^2}{2m_n^*}; E = E_V - \frac{\hbar^2 k^2}{2m_p^*}, \quad (11.14)$$

So that

$$h\nu - E_g = \frac{\hbar^2 k^2}{2m_r^*}, \quad (11.15)$$

where $E_g = E_C - E_V$ - is the latitude of the forbidding zone, and m_r^* - electron, which is the cited mass of the pit, which is determined by the following attitude:

$$\frac{1}{m_r^*} = \frac{1}{m_n^*} + \frac{1}{m_p^*}. \quad (11.16)$$

According to (11.6), α the absorption coefficient $\nu(h)$ is inversely proportional to the average free travel path of the photon l_f , which in turn is determined by the speed of the photon ν_f and its average free path time τ_f . So,

$$\alpha(h\nu) = 1/l_f = 1/\nu_f \tau_f = g(h\nu)\pi/c. \quad (11.17)$$

where $g(h\nu) = 1/\tau_f$ is the probability of a photon absorption in a unit of time; $\nu_f = c/\pi$ - photon velocity in the medium; π - refractive index of a substance (the actual part of the indicator); C - The velocity of light in a vacuum.

The number of light quanta with absorbed energy in the unit of ν time in the range from h to $h(\nu+d\nu)$ is equal to the probability of the electron passing $P(h)\nu$ and $|E|$ The voltage in the energy range is proportional to the number of quantum states in the stress zone:

$$g(h\nu)d(h\nu) = 2P(h\nu)N(E)|dE|. \quad (11.18)$$

The coefficient of 2 here takes into account the two possible directions of light polarity. According to (4.11), the density of quantum states in the stress band is as follows:

$$N(E) = \frac{(2m_p^*)^{3/2}}{2\pi^2\hbar^2}(E_V - E). \quad (11.19)$$

Using (11.14) and (11.19), we get:

$$g(h\nu)d(h\nu) = \frac{2}{\pi^2}P(h\nu)k^2dk. \quad (11.20)$$

But (11.15) follows from the expression:

$$d(h\nu) = \frac{\hbar^2}{m_r^*}kdk, \quad (11.21)$$

here

$$k = \frac{(2m_r^*)^{1/2}}{\hbar}(h\nu - E_g)^{1/2} \quad (11.22)$$

So that

$$g(h\nu) = \frac{(2m_r^*)^{3/2}}{\pi^2\hbar^3}P(h\nu)(h\nu - E_g)^{1/2}. \quad (11.23)$$

With the preservation of the quasi-impulse (11.11) for the directly permissible transitions, the probability of the transition does not depend on the photon energy of $P(h\nu)$. Given this, the private absorption coefficient for direct allowable jumps based on (11.17) and (11.23) (see Fig. 11.2) is expressed by the following formula:

$$\left\{ \begin{array}{l} \alpha = A(h\nu - E_g)^{1/2}, h\nu > E_g; \\ \alpha = 0, h\nu \leq E_g, \end{array} \right\} \quad (11.24)$$

here

$$A = \frac{\pi e^2 [2m_n^* m_p^* / (m_n^* + m_p^*)]^{3/2}}{h^2 c m_n^* \varepsilon_r \pi}. \quad (11.25)$$

(11.24) The relation $h\nu - E_g$ applies within a range of limited variability. Fig. 11.4 shows that α^2 are linearly related to g_a at a given interval $h\nu$. Continuing this

line until it intersects with the abscissa axis allows you to determine the width of the restricted zone E_g for the permissible crossings directly .



Fig. 11.4. The energy dependence of adsorption coefficient squared for straight allowable jumps

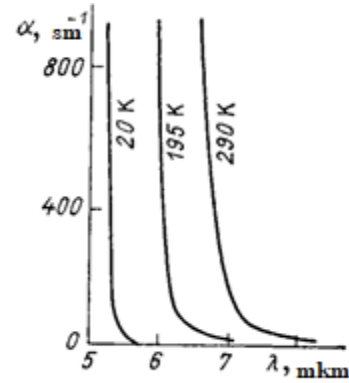


Fig. 11.5. Absorption spectrum of surrogate antimonoxide at different temperatures

It follows from (11.24) that there should be no absorption of energetic quanta that is less than the width of the forbidden zone for straight transitions. Therefore, the private absorption edge on the side of the long waves (or small energies) must be very sharp. Indeed, as shown in Fig. 11.5, there is a sharp increase in absorption in a very pure antimonide indium monocrystal with the right zones.

Transitions from the stress zone in the s-state to the conductivity zone based on the atomic p-function would be direct, but "forbidden" (see Fig. 2.11). For all $k \neq 0$, the probability of passing in the forbidden transitions is proportional to $P(h\nu)k^2$.

As a result of this, the probability of a pass for straight transitions that is prohibited depends on the energy as follows:

$$P(h\nu) \approx k^2 = \text{const}(h\nu - E_g). \quad (11.26)$$

Therefore, the absorption coefficient for properly prohibited inter-zone crossings is expressed by the following attitude:

$$\alpha'(h\nu) = A'(h\nu - E_g)^{3/2}. \quad (11.27)$$

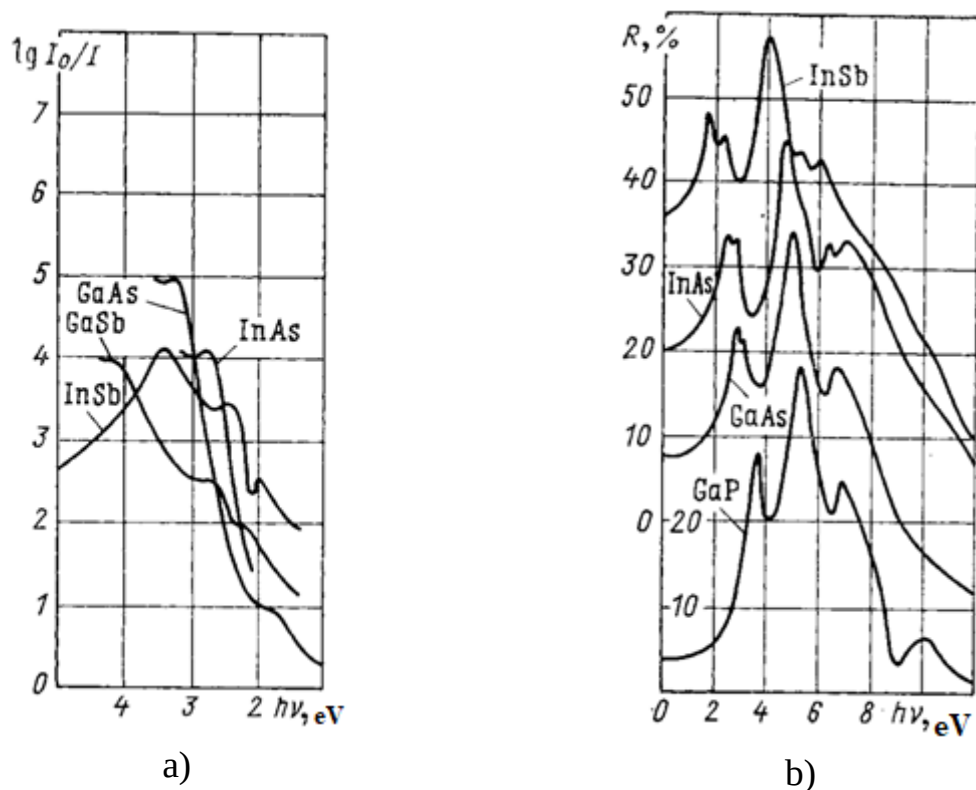


Fig. 11.6. Absorption spectra of some $A^{III}B^V$ semiconductors at 80 K (a) and return spectra at 300 K (b)

The conductivity zone of semiconductors usually contains several small bands (see Fig. 2.21). Therefore, when a photon whose energy is greater than the width of the bandwidth is absorbed, direct transitions from the stress band to the conductivity band can occur at all points in the quasi-impulse space where the selection rule is satisfied. In this case, the absorption spectrum returning from the complex structure of the conductivity zone becomes complex. For such researches, thin films of semiconductors are used, since the absorption coefficient in photon energies of $E > E_g$ is $\propto$ very large. As an example, Fig. 11.6-a shows the absorption spectra of some $A^{III}B^V$ -type semiconductor thin films at 80 K.

At higher energies, the right transitions also occur on the return spectra. In them, as can be seen from Fig. 11.6-b, a number of peaks are observed, reflecting the complex structure of the permeability zone. The high intensity of these peaks is due to the high density of the initial and last cases.

11.3. Private Takeover in Improper Transitions

In some cases, there are factors of "softening" the rules of selection, as a result of which vertical transitions between non-straight zones are also accepted (Transition 2 in Fig. 11.2b). However, the probability of such transitions is much smaller than the probability of straight transitions (Fig. 1 in Fig. 11.2-b). In contrast to straight transitions, the law of conservation of quasi-impulse in false transitions

is provided by its interaction with the phonon, as shown in Fig. 11.7-a. Therefore, non-biased or non-vertical jumps occur by emission or absorption of phonones. These jumps determine the absorption that is located on the long-wave side of the private absorption boundary caused by the correct transitions. The right and false transitions are very clearly manifested in the absorption spectrum of germanium and silicon. According to Fig. 11.8, the basic location values of the phonon energy, which correspond to the false and correct optical transitions at room temperature, are 0.64 and 0.81 eV, respectively.

In a semiconductor with fault transition zones, a fault transition can occur from any band state of the stress band to any free state of the conduction band. The absorption coefficient for the false transitions between fault zones that occur in conjunction with the absorption or emission of a single phonon can be calculated as follows. We consider a false transition as a transition that occurs through a series of intermediate "vertical" states with very short residence times. In this case, the law of conservation of energy is fulfilled only for the whole transition, when the electron passes into a virtual state, its quasi-impulse is conserved.

For simplicity's sake, we assume that the false transition between the fault zones occurs from the initial state $k_1=0$ to the final state $k_2=k_{min}$ (Transition 2 in Figure 11.7-b). There are two possibilities for such a transition, which occurs with the output or absorption of the phono. An electron located near the top of the stress zone where $k_1=0$ is excited by light and passes into the conductivity zone, where it occupies a position with the same wave vector (Transition 1 in Fig 11.7-b). As a result of such a transition in the stress zone, a hole with $k_1=0$ remains.

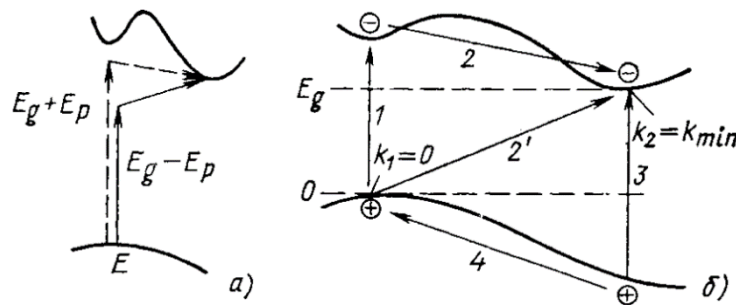


Fig. 11.7. Fault Zones Exist in Semiconductor Fault Transitions

At the same time , it has a higher energy than the energy corresponding to the lower part of the electron conductivity zone where the electric conductivity is $k_1=0$, so that the minimum energy absorbed or ejected the phonon over a very short time interval passes into the state near $k_2=k_{min}$ (Fig. 2 in Fig. 11.7-b).

It can also be a jump as follows. Having absorbed a quantum of light, the electron vertically transitions from a state lying deep in the stress zone to a state

where $k_2=k_{min}$ (11.7-b Fig. 3) in the conductor zone. As a result, a deep-lying hole is left in the stress zone, which passes into a position near the top of the stress zone due to phonon emission or absorption (11,7 b, Fig. 4).

We calculate the absorption coefficient for the transitions when the electron makes the vertical first jump and then the second jump. Let the initial energy of the electron in the stress zone be E , let the energy of the electrons in the last state in the conduction zone be E' . The energy of the phonon involved in optical transitions is determined by E_p . If the energy calculation begins at the bottom of the stress zone, then the energy of the quantum of light (Fig. 11.7 b) required to perform the reverse transition to the electron in the case of absorption of the phonon is determined by the following equation:

$$h\nu_\alpha = E' - E - E_p, \quad (11.28)$$

The energy of the quantum of light required to perform the reverse transition to the electron in the case of photon emission

$$h\nu_e = E' - E + E_p \quad (11.29)$$

So so.

The absorption coefficient α is proportional to the integral over all possible pairs of the densities of the initial $N(E)$ and the final states $N(E/h\nu \pm E_p)$ by energy, which means that,

$$\alpha = \int_0^{E_m} N_p N(E) N(E') dE. \quad (11.30)$$

Taking into account the received energy (Fig. 11.7), the density of the initial states at energy E (in the tensile zone) is as follows:

$$N(E) = \frac{(2m_p^*)^{3/2}}{2\pi^2 \hbar^3} |E|^{1/2}, \quad (11.31)$$

and based on (11.28) and (11.29), the density of the states in E' energy (in the permeability zone) is equal to:

$$N(E') = \frac{(2m_n^*)^{3/2}}{2\pi^2 \hbar^3} (E' - E_g)^{1/2} = \frac{(2m_n^*)^{3/2}}{2\pi^2 \hbar^3} (h\nu - E_g \pm E_p + E)^{1/2}. \quad (11.32)$$

If the non-equilibrium phonon distribution function is not much different from the balanced (n), then it can be assumed that the number of phonons based on (3.75) is determined by the following formula:

$$N_p = \langle n \rangle = \frac{1}{e^{E_p/kT} - 1}. \quad (11.33)$$

$$E_m = -(h\nu - E_g \pm E_p),$$

$$\alpha(h\nu) = A' N_p \int_0^{-(h\nu - E_g \pm E_p)} |E|^{1/2} (h\nu - E_g \pm E_p + E)^{1/2} dE = AN_p (h\nu - E_g \pm E_p)^2, \quad (11.34)$$

Here A-is some kind of constant magnitude.

If we put (11.33) in (11.34), the absorption coefficient for the transitions that occur with the absorption of the background takes the following form:

$$\alpha_\alpha(h\nu) = \frac{A(h\nu - E_g - E_p)^2}{e^{E_p/kT} - 1}, \quad h\nu > E_g - E_p; \quad (11.35)$$

$$\alpha_\alpha = 0, \quad h\nu \leq E_g + E_p.$$

The probability of phonon emission is proportional to N_p+1 , hence the absorption coefficient in phonon emission:

$$\alpha_e(h\nu) = \frac{A(h\nu - E_g + E_p)^2}{1 - e^{-E_p/kT}}, \quad h\nu > E_g + E_p. \quad (11.36)$$

$h\nu > E_g + E_p$ Since light can be absorbed by both the absorption of phonons and the emission of phonons, the sum of the absorption coefficient for non-linear transitions will be:

$$\alpha(h\nu) = \alpha_\alpha(h\nu) + \alpha_e(h\nu) = A \left[\frac{(h\nu - E_g + E_p)^2}{e^{-E_p/kT} - 1} + \frac{A(h\nu - E_g - E_p)^2}{1 - e^{-E_p/kT}} \right], \quad h\nu > E_g + E_p. \quad (11.37)$$

(11.37) attitude can be summed up as follows:

$$E_g - E_p < h\nu \leq E_g + E_p \text{ For } \alpha = A \frac{(h\nu - E_g + E_p)^2}{e^{-E_p/kT} - 1}, \quad (11.38)$$

$$h\nu \leq E_g - E_p \text{ for } \alpha = 0.$$

The graph of the dependence of $\alpha^{1/2} h\nu$ for germanium and silicon is shown in Fig. 11.8-9 below, with two straight linear plots.

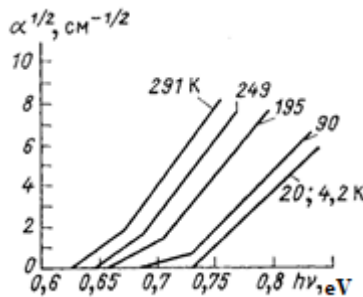


Fig. 11.8. Dependence $\alpha^{1/2}$ on $h\nu$ for German

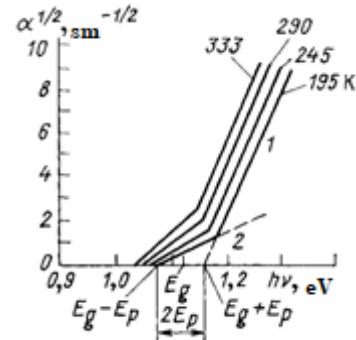


Fig. 11.9. Dependence of $\alpha^{1/2} h\nu$ for silicon

Straight line 1 corresponds to the intersecting relationship with the abscissa axis at the $\alpha_e^{1/2}$ point $E_g + E_p$. Straight line 2 describes the $\alpha_\alpha^{1/2}$ connection and intersects with the abscissa axis at the point $E_g - E_p$. The $\alpha_\alpha^{1/2}$ cross-section length $\alpha_e^{1/2}$

between the points of intersection corresponding to the same temperature and ν on the h -axis of the lines is $2E_{p\nu}$. In the middle of this cross-section is the point $h\nu = E_g$. At the same time, as shown in Figure 11.9, when the study temperature decreases, the $\alpha^{1/2}$ dependency graph $h\nu$ shifts towards larger values of h , reflecting the temperature dependence of the width of the forbidden zone, and the $\alpha^{1/2}$ slope of the straight line tends to zero. This is because at low temperatures very few phonons are excited, and therefore there cannot be an optical transition by absorption of the phonon.

If a correct crossing is prohibited at $k_I=0$, then the incorrectly prohibited inter-zone crossings for this case are characterized by a absorption coefficient, according to (11.26):

$$E_g - E_p < h\nu \leq E_g + E_p \text{ uchun } \alpha \sim \frac{(h\nu - E_g + E_p)^3}{e^{-E_p/kT} - 1}. \quad (11.39)$$

The absorption coefficient for direct permissible inter-zone crossings α generally reaches 10^4 - 10^5 cm^{-1} and for non-permissible inter-zone crossings is $\alpha = 10^{-1}$ - 10^3 sm^{-1} (see Fig. 11.8).

Only one type of phonon was considered when considering inappropriate cross-zone transitions. In fact, a crystal can contain many, many different phonons. Optical transitions involve only phonons that provide the desired change in the electron quasipulse. Such phonons are typically longitudinal and transverse acoustic phonones. Table 11.1 shows the values of E_p phonon energy for germanium and silicon, as well as $h\theta$ equivalent temperature $\theta = E_p/k$.

Table 11.1. Debay temperatures for the background energy and different networks of vibrations in germanium and silicon

Semiconductor	Type of oscillations	Phonon energy $E_p = k\theta$, eV	Equivalent temperature θ , K
Germanium	Longitudinal acoustic LA	0.0275	320
	Acoustic Guitar	0.0077	90
	Longitudinal optical TO	0.034	400
Silicon	Longitudinal acoustic LA	0.016	230
	Acoustic Guitar	0.055	700
	Longitudinal optical TO	0.119	1370
	Longitudinal optical LO	0.083	1010

11.4. Absorption in Highly Alloyed and Amorphous Semiconductors

With an increase in the concentration of free electrons in the indium antimonide, which has a properly banned zone, the edge of the private absorption

zone shifts towards shorter wavelengths. For example, for an n-type material with an electron concentration of $5 \cdot 10^{18} \text{ cm}^{-3}$, the adsorption limit at room temperature is $3,5 \text{ mkm}$, whereas for an unalloyed semiconductor, the long-wave absorption limit at room temperature is $7,1 \text{ mkm}$ (Fig. 11.10 A).

This is due to the fact that the effective mass of electrons in indium antimonide is characterized by ($m_n^* = 0.013m_0$) a large curvature due $\frac{1}{m^*} = \frac{1}{\hbar} \frac{\partial^2 E}{\partial k^2}$ to the small of the conduction zone, so that the density of the states in it is small (Fig. 11.10 b).

As a result, with an increase in the concentration of free electrons, the spinning of the electron gas occurs rapidly, and the Fermi level rises from the lower part of the F_n conduction zone to $\xi_n = F_n - E_C$. Since cases below the Fermi level have already been filled, there is no breakthrough associated with the transition to these cases. In this case, $h\nu_{min}$ (the correct transition in Fig. 11.10 b) would be greater than the width of the forbidding zone, and therefore the absorption edge ξ_n should shift to higher energies of approximately.

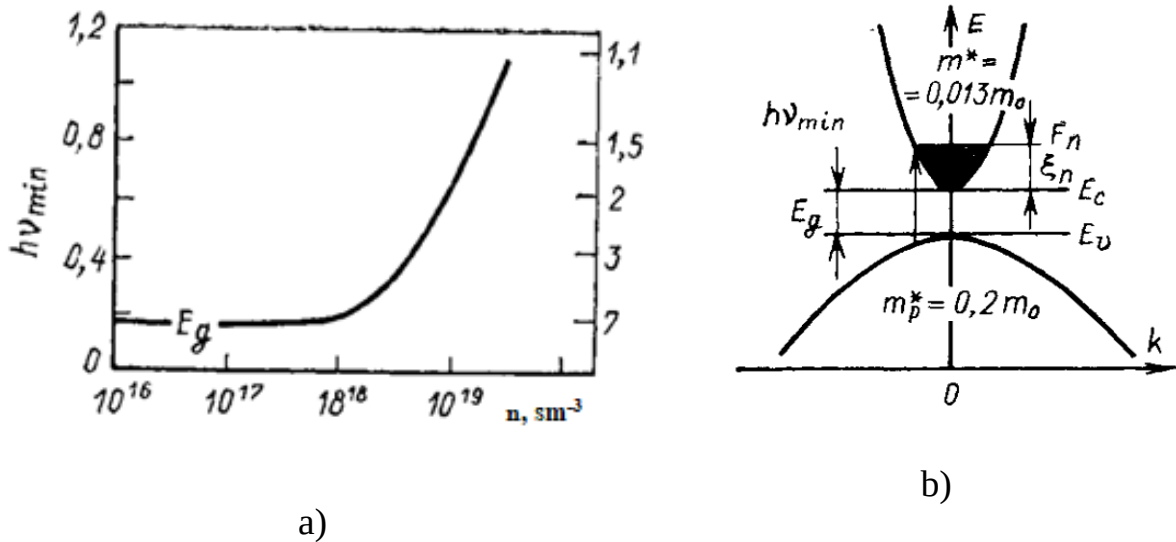


Fig. 11.10. Dependence of the optical energy range in n-type InSbv on the electron concentration (a) and the energy on the quasiimpulse (b). sm^{-3}

11.5. Influence of External Factors on Private Semiconductor Disposal

Many experiments conducted on various semiconductors show that external influences, such as temperature, pressure, electric and magnetic fields, have significant effects on their individual absorption spectrum. Knowing their effect on the energy levels of a semiconductor is necessary both to solve theoretical problems related to considering the state of electrons in a crystal and to solve problems of practical relevance related to the manufacture and use of semiconductor devices.

11.5.1. The influence of temperature

As the temperature increases in the semiconductor, the vibrations of the atoms relative to the equilibrium state increase, and the lattice expands. For some semiconductors, a reduction in the width of the bandwidth with an increase in lattice constant, according to Fig. 2.7, is observed along with the displacement of the main absorption band towards the longer wavelengths (see Fig.11.5). For many semiconductors in temperature ranges well below the Debay temperature, the temperature dependence of the bandwidth of the bandwidth is characterized by an empirical relationship in the following form:

$$E_g(T) = E_g(0) - \frac{\alpha T^2}{T + \beta}, \quad (11.40)$$

where $E_g(0)$ – 0 K the width of the prohibited zone at; α and β are constant values, and their values for some semiconductors are shown in Table 11.2.

Table 11.2. E_g Value and α β Parameters

Yarimo'tkazgich Has	Type of inter- zone	$E_g(0)$, eV	$\alpha \cdot 10^4$	B
Yes	Egi	1,1557	7,021	1108
GE	Egi	0,7412	4,561	210
GE	Egd	0,8893	6,842	398
6H-SiS	Egi	3,024	-0,3055	-311
GaAAs	Egd	1,5216	8,871	572
InP	Egd	1,4206	4,906	327
InAs	Egd	0,426	3,158	93

Here 's the E_{gi} Mispasses; E_{gd} - straight transitions.

At temperatures much higher than the Debay temperature, the dependence of the forbidden zone width is linear:

$$E_g(T) = E_g(0) - \alpha T. \quad (11.41)$$

11.5.2. Effects of Pressure

When a crystal is compressed in all directions, called hydrostatic pressure, the distance between its atoms decreases. However, according to Fig. 2.7, the reduction in interatomic distance should result in an increase in the width of the prohibited zone.

The increase in E_g with increased pressure is observed in most semiconductors that have proper bandwidths. With the increase in hydrostatic pressure in such semiconductors, the private absorption limit shifts towards a shorter wavelength.

The pressure dependence of E_g in non-bandwidth semiconductors can be complicated. For example, in Germanium, as shown in Fig. 11.15, E_g increases with increasing pressure, while for silicon, on the contrary, it decreases (Fig. 11.16). With an increase in hydrostatic pressure in gallium arsenide, the minimum of the conductivity zone for straight transitions moves upwards faster than the minimum for false transitions. Therefore, when the pressure increases, it can become a semiconductor with false transitions.

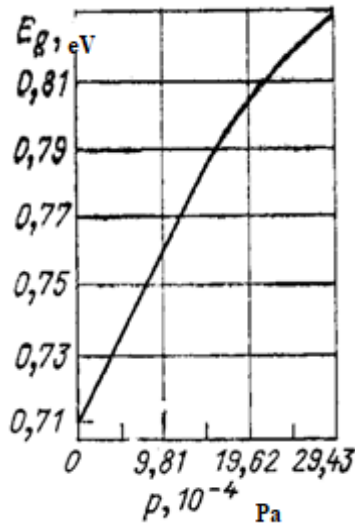


Fig. 11.15. Pressure dependence of the width of the German exclusion zone at 350 K

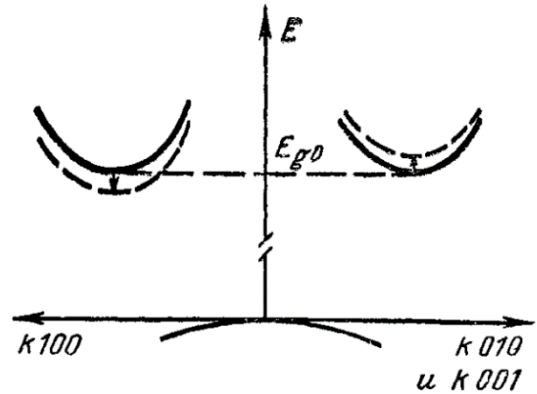


Fig. 11.16. Longitudinal and transverse displacement of the silicon permeability zone under the influence of single-axis compression in the direction of [100]

One-way compression or elongation of the crystal causes a shift of the state of different energy minimums in the conductivity zone. For example, the compression of a silicon crystal along its axis [100] reduces the width of the forbidding band in the direction of compression, but due to its simultaneous stretching in the transverse direction, the minimums of the conductivity zone located in the plane perpendicular to the direction of compression shift into the high-energy region (Fig. 11.16). In n-type germanium, the greatest impact is observed under compression in the direction of [110]. The change in the energy structure of the crystal causes a redistribution of electrons, where the electrons pass to smaller energy states as a result of shear deformation and affect the value of the elastic constants, i.e. the coefficient of pesoresistance.

11.5.3. The Effect of a Strong Electric Field

As defined in paragraph 2.10, the energy band of a semiconductor bends in the external electric field. In this case, the electrons of the stress zone can be tunneled into the conduction zone.

According to Fig. 11.17, according to (6.127), the height of the triangular barrier at the bend of the energy zones E_g is equal to a and the thickness is d :

$$d = E_g / eE. \quad (11.42)$$

As can be seen from the expression, as the external electric field strength increases E , the barrier thickness decreases, and as a result, the probability of the electron tunneling process increases.

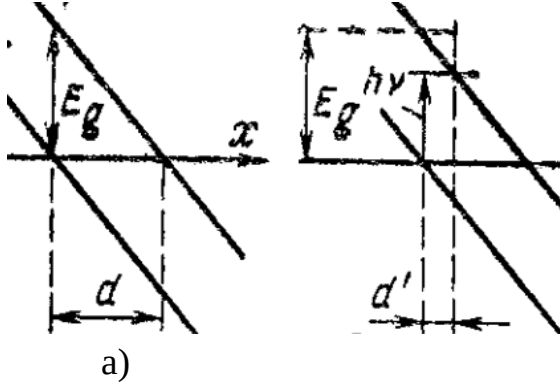


Fig. 11.17. Electron tunneling: a - without changing energy; b - with the decay of the photon

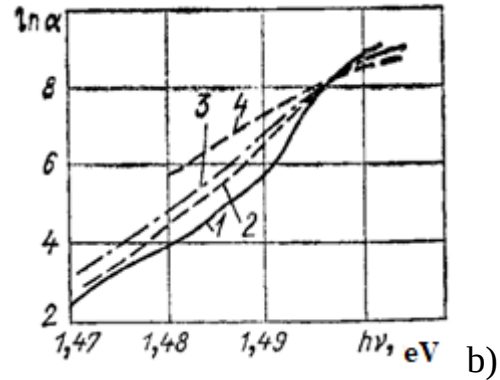


Fig. 11.18. GaAs absorption limit at different values of electric field strength: 1 - zero; 2 - 6 kV/cm; 3 - 10 kV/cm; 4 - 30 kV/cm

As can be seen from Fig. 11.17-b, when an external electric field is present, there is an absorption of light quanta with an energy of less $h\nu$ than the width of the semiconductor's forbidden zone, which is equivalent to a decrease in the barrier thickness to a value d' . In this state

$$d' = (E_g - h\nu) / eE \quad (11.43)$$

And the likelihood of tunnel crossing increases.

The process of tunneling electrons involving an external electric field, which occurs with the absorption of photons, is commonly referred to as the Frantz-Keldisch effect. In the corresponding semiconductor, it manifests itself as a shift of the absorption edge to lower energies (Fig. 11.18).

11.5.4. The Effect of a Magnetic Field

Let us first consider the effect of a magnetic field on the energy spectrum of a crystal. This problem can be solved on the basis of the Schrödinger equation, which states that the Hamiltonian should determine the state of an electron in a magnetic field.

From the theory of electromagnetic fields, it is known that the magnetic field induction vector B is with vector potential $A(x, y, z)$

$$\vec{B} = \text{rot} \vec{A} \quad (11.44)$$

Connected through attitude.

In this case, an electron with charge $-e$ and impulse p acquires a generalized impulse when passing through the field of impact of the field determined by the vector potential $A(x, y, z)$.

$$\vec{P} = \vec{p} + e\vec{A} \quad (11.45)$$

In this case, since it is kinetic energy $T = \frac{p^2}{2m_0}$, the Hamiltonian function for the electron is written as follows.

$$H = T + U(x, y, z) = \frac{1}{2m_0} (\vec{p} + e\vec{A})^2 + U(x, y, z), \quad (11.46)$$

But $U(x, y, z)$ - energy potential.

For accuracy, let's assume that a homogeneous magnetic field is directed along the z -axis, namely:

$$B_x = B_y = 0; B_z = B. \quad (11.47)$$

In this

$$\text{rot} \vec{A} = \begin{vmatrix} i & j & k \\ \frac{\partial}{\partial x} & \frac{\partial}{\partial y} & \frac{\partial}{\partial z} \\ A_x & A_y & A_z \end{vmatrix} = i \left(\frac{\partial A_z}{\partial y} - \frac{\partial A_y}{\partial z} \right) + j \left(\frac{\partial A_x}{\partial z} - \frac{\partial A_z}{\partial x} \right) + k \left(\frac{\partial A_y}{\partial x} - \frac{\partial A_x}{\partial y} \right)$$

Based on (11.47):

$$\left. \begin{aligned} B_x &= \frac{\partial A_z}{\partial y} - \frac{\partial A_y}{\partial z} = 0; \\ B_y &= \frac{\partial A_x}{\partial z} - \frac{\partial A_z}{\partial x} = 0; \\ B_z &= \frac{\partial A_y}{\partial x} - \frac{\partial A_x}{\partial y} = B. \end{aligned} \right\} B_z = B. \quad (11.48)$$

(11.48) links

$$A_x = A_z = 0; A_y = Bx. \quad (11.49)$$

satisfies us.

Given (11.49), the Hamiltonian function (11.46) is defined in relation to:

$$H = \frac{1}{2m_0} [p_x^2 + (p_y^2 + eA_y)^2 + p_z^2] + U(x, y, z). \quad (11.50)$$

Using the approximation of the effective mass (see paragraph 2.14), we write the Schrödinger equation $\hat{H}\Psi = E\Psi$. From the Hamilton function to the Hamilton

operator (11.50), we pass through the following $p_x - \frac{\hbar}{i} \frac{\partial}{\partial x} ga$, $p_y - \frac{\hbar}{i} \frac{\partial}{\partial y} ga$ and $p_z - \frac{\hbar}{i} \frac{\partial}{\partial z} ga$ substitutions. As a result, we get:

$$-\frac{\hbar^2}{2m^*} \left[\frac{\partial^2 \Psi}{\partial x^2} + \left(\frac{\partial}{\partial y} + \frac{ieBx}{\hbar} \right)^2 \Psi + \frac{\partial^2 \Psi}{\partial z^2} \right] = E\Psi. \quad (11.51)$$

Control Questions

1. Describe the return spectrum
2. Explain the Absorption Spectrum
3. What do you mean by private breakthrough in the right transitions?
4. Explain Private Swallowing in False Passes
5. How does absorption occur in highly alloyed and amorphous semiconductors?
6. What will be the effect of temperature on the private absorption of semiconductors?
7. Explain the effect of pressure on private absorption of semiconductors
8. Describe the effect of a strong electric field on the private absorption of semiconductors
9. What is the effect of the magnetic field on the private absorption of semiconductors?

Chapter XII.

LUMINESCENCE IN SEMICONDUCTORS

12.1. Types of luminescence

In semiconductors, in the excited state, the opposite process to absorption, the output-luminescence of electromagnetic radiation, can be observed. This state of the substance can be acquired in different ways. In photo luminescence, radiation occurs due to absorbed light energy. The flash that occurs when a sample is attacked with electrons is called cathodoluminescence. Luminescence can occur when a substance is exposed to an electric field (electroluminescence) or mechanical effects (triboluminescence), X-ray radiation (X-ray radiation), and chemical and biological processes (chemical and bioluminescence).

Basically, there are two types of luminescence in solids: monomolecular and recombination.

12.2. Monomolecular irradiation of solids

Solid-state monomolecular radiation occurs when light is absorbed and emitted at the center of ingress. The luminescence spectra of solids, i.e., the dependence of the intensity of light on the energy of the emitted photon or wavelength, differ from the corresponding spectra of isolated ingress atoms in two main respects. Firstly, the radiation of solids is usually observed in the form of wide streaks, and secondly, their luminescence spectrum is usually shifted towards long wavelengths relative to the absorption spectrum, that is, the energy emitted is less than the absorbed energy. The difference between these two energies is called the "Frank-Condon shift." This peculiarity of the optical properties of solids is due to the interaction of the center of radiation with the crystal lattice area of their appearance. When the center interacts with its immediate surroundings, its potential energy consists of Coulomb gravity, alternation interaction, and van der Waals gravity. As a result of these interaction forces, as shown in Fig. 12.1, the levels of input atoms in the main U_0 and excited U_q states in the crystal lattice are quasmolecular.

In such a situation, the light can be absorbed by the inlet center, as a result of which the electron U_0 passes from the main state to the state where U_q is concentrated. The spontaneous inverse transition of the electron from the centrifugal state to the ground surface is accompanied by radiation. If we consider the occupancy rates of electrons in the primary and excited states, then the corresponding set of photons forms absorption and ejection lines, which are illustrated on the energy scale on the axis of coordinates in Fig. 12.1. The duration

of such luminescence of a solid is determined by the residence time of the electron in the central excited state. It is immutable, independent of external factors, and, in particular, for the dipole radiation of atoms, it is about 10^{-8} s.

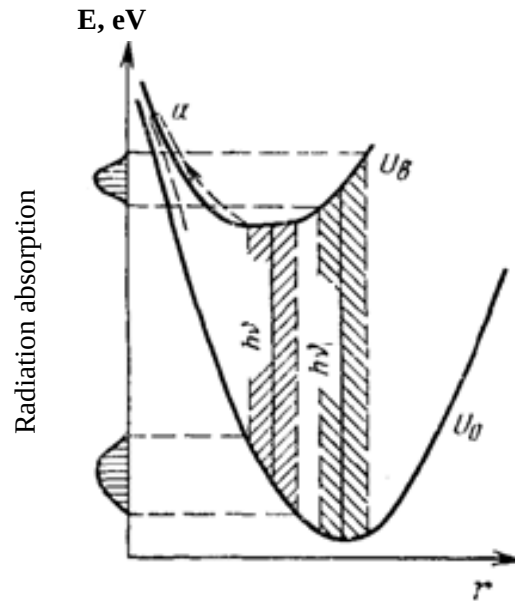


Fig. 12.1. Diagram of energy levels of a solid-state ingress center

12.3. Recombinatory radiation of semiconductors at fundamental transitions

When electron-pit pairs are formed by semiconductor entitation, recombination radiation can occur from a direct recombination of free electrons and free swells, annihilation of excitons, as well as recombination of free electrons and free sleds via recombination traps that act as light centers. First, we will consider fundamental transitions that are equal to or close to the width of the zone in which the energy change is prohibited.

12.3.1. Straight transitions "permeability zone - tensile zone"

The absorption of light in a straight-bandwidth semiconductor occurs in conjunction with the formation of a free electron and a free pit where the wave vectors are identical ($k = k'$). The free charge carriers formed after the process of absorption of light undergo scattering, due to which, during relaxation, they fall to the bottom of the electron conductivity zone and rise to the ceiling of the cavity stress zone. This state between the free charge carrier and the crystal grid is established in about 10^{-10} - 10^{-12} s. If the wave vectors of the electron and the hive are the same after the relaxation process, the correct recombination of the electron and the free hive accompanied by photon radiation is highly likely to occur (Fig. 12.2).

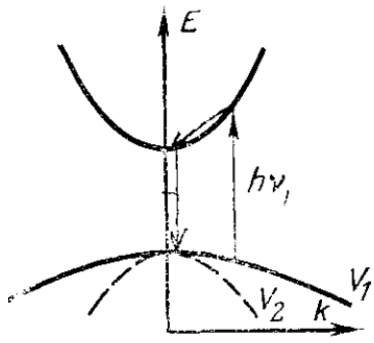


Fig. 12.2. Straight beam jumps

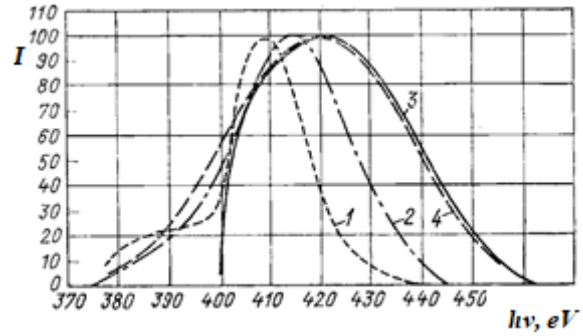


Fig. 12.3. Radiation of n-type indium antimonide at 77 K 1-2.3·10¹⁶ cm⁻³; 2-9·10¹⁶ cm⁻³; 3 - 1.8·10¹⁷ cm⁻³ (theory); 4-1.8·10¹⁷ cm⁻³

At this point, the radiation spectrum is characterized by the following expression, just like that of corresponding absorption processes:

$$I(h\nu) = B(h\nu - E_g)^{1/2} \quad (12.1)$$

Here's a B -variable magnitude.

This formula shows that during the recombination of free charge carriers, the radiation must have a threshold on the low-energy side at $h\nu = E_g$ (Fig. 12.3, Curve 3). With the increase in temperature, as well as with an increase in excitation energy, the upper cases in the superconductivity zone are filled. This results in radiation at high photon energies that produce a high-energy "tail" in the luminescence spectrum.

As an example of direct radiation interband recombination, Fig. 12.3 shows the photoluminescence spectrum of n-type InAs with varying degrees of exposure. As the degree of legering increases, the radiation maximum and the displacement of its high-energy edge towards the high-photon energy side is related to the displacement of the Fermi state into the conductivity zone.

12.3.2. False transitions "permeability zone - tensile zone"

The absorption of light in a non-randomly banded band semiconductor occurs when the phonon is absorbed or emitted. In such a semiconductor, the phonons that ensure the conservation of the quasi-impulse must also be involved in the recombination of a free electron and a free cell (Fig. 12.4).

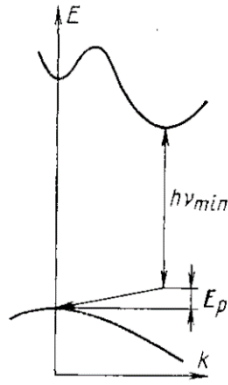


Fig. 12.4. False radiation transitions

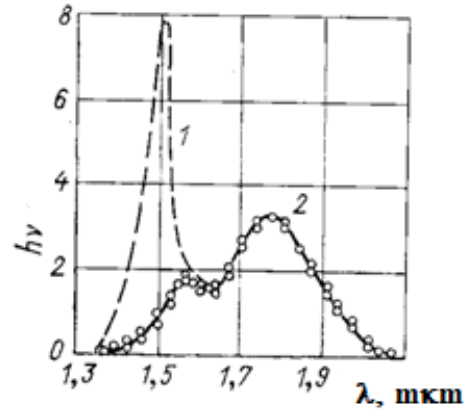


Fig. 12.5. Germanium radiation spectrum (sample thickness $1.3 \cdot 10^{-3} \text{ cm}^{-3}$; 2 - curve experimental; Curve 1 is obtained taking into account reabsorption

The process with the highest probability is phonon emission. The process of absorption of a phonon is irrelevant with respect to the process of phonon ejection, since its probability is less than the probability of phonon emission is $N_p(N_p+1)$ times. The optical transition that occurs with the ejection of the phonon has an energy equal to $h\nu_{min} = E_g - E_p$, less than the width of the forbidding zone. In this case, similar to the absorption coefficient, the radiation spectrum is characterized by the following expression:

$$I(h\nu) = B(h\nu - E_g + E_p)^2 \quad (12.2)$$

where E_p is the energy of the phonon involved in recombination.

If forward and reverse transitions are observed in a semiconductor, two luminescence outliers are observed in its radiation spectrum, as in germanium, for example, (Fig. 12.5). Long wave radiation streaks are determined by non-emission recombination. Short-wave luminescence emissions are associated with radiating electron transitions from the straight zone located at an altitude of 0.15 eV from the bottom of the lowest zone. Curve 2 in this figure is experimental, and curve 1 is obtained from curve 2 by making adjustments for absorption in the sample.

12.3.3. Highly Alloyed Semiconductor

In a strongly alloyed semiconductor, the Fermi level lies within the permissible zone (in the conductivity zone for n-tip and in the stress zone for p-tip). Fig. 12.6 shows the energy diagram of the semiconductor and shows the formation of electron concentrations in the conduction zone due to excitation.

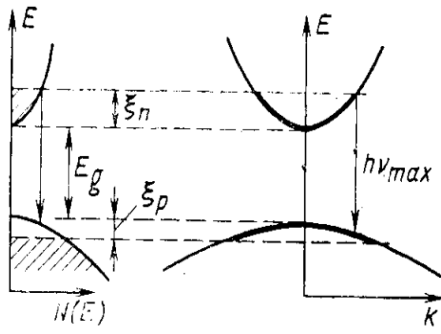


Fig. 12.6. Correct in sober condition
Jumps

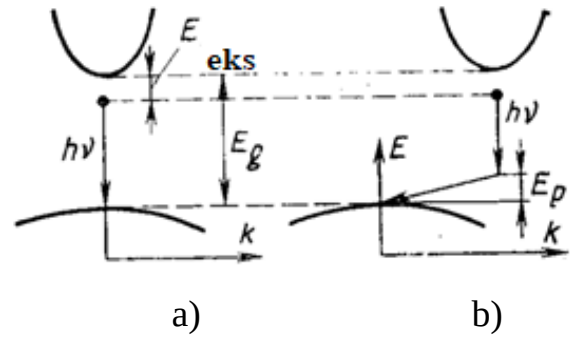


Fig. 12.7. Exciton recombination
A-To'ji'ri; B – Noto'ji'ri O'Tischler

In semiconductors with a straight bandwidth for straight jumps, the intensity of radiation should increase as follows:

$$I(h\nu) = B(h\nu - E_g)^{1/2} \quad (12.3)$$

In non-transient transitions, due to electron-electron scattering or "electron-input" scattering processes, the law of conservation of a quasi-impulse is performed, all occupied states in the conduction zone can cause radiant transitions to all empty states of the stress band. In this case, the same energy $h\nu$ can correspond to the transitions between the different upper and lower states, so the radiation spectrum is characterized by the following expression:

$$I(h\nu) = B(h\nu - E_g)^2. \quad (12.4)$$

12.3.4. Exciton recombination

In a sufficiently pure semiconductor, when excited by light whose energy is close to the width of the forbidden zone, a free electron and a free spark bound by the exciton-Coulomb gravitational force can occur. When such charge carriers are recombined, that is, by exciton annihilation, a narrow spectral line is radiated. In a semiconductor with straight-bandwidth permissible transitions (Figure 12.7, a), the energy of the exciton radiation is:

$$h\nu = E_g - E_{eks}, \quad (12.5)$$

This is where the energy of the Ecclesiastes is located.

According to E_{eks} (11.71), it is possible to have excitation states where the ionization energy is n^2 times less than the ionization energy of the primary state, corresponding to $n = 1$. Therefore, the emissivity of the free excitation may consist of a series of narrow lines.

In a semiconductor with a falsely forbidden zone (Fig. 12.7-b), i.e., the law of conservation of a quasi-impulse, the energy of a radiated photon is realized by the emission of an optical phonon with an energy of E_p :

$$h\nu = E_g - E_{eks} - E_p. \quad (12.6)$$

In semiconductors with correctly fixed transitions, free exciton radiation can also occur by the emission of one or more phonons (Fig. 12.8), but the likelihood of such a transition decreases and the corresponding glow of radiation becomes weaker.

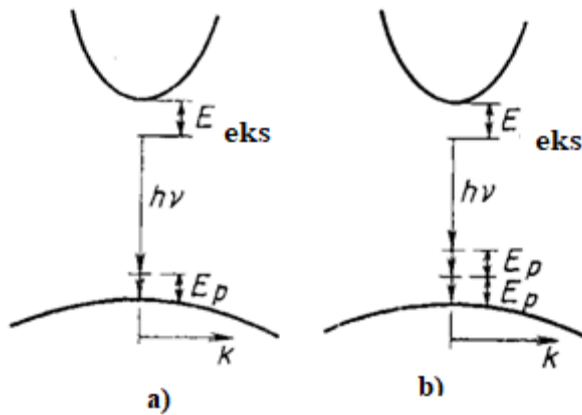


Fig. 12.8. The correct exciton recombination is: a - with the output of a single optical phonon; b - with output of two optical backgrounds

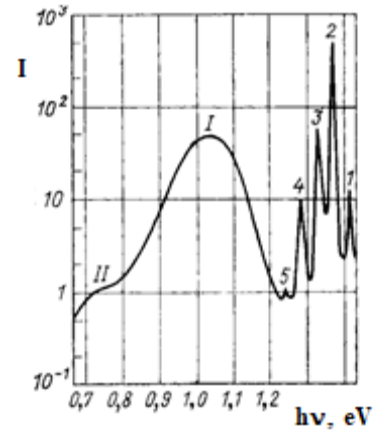


Fig. 12.9. Photoluminescence spectrum of InP at 6 K

In the presence of penetrations, connected excitons can form. Their radiation spectrum consists of narrow bands that are smaller than the recombination energy of the free exciton. Fig. 12.9 shows the photoluminescence spectrum of much purer indium phosphide as an example. The radiation peak denoted by the number 1 is associated with the recombination of the free exciton, while lines 2, 3, 4, and 5 are associated with the recombination of the bound exciton, corresponding to the output of the longitudinal optical phonons 0, 1, 2, and 3 with an energy of 43 meV. Radiation lines I and II occur as a result of non-excitation processes.

12.4. Recombination radiation at the transition between zones and access surfaces

In semiconductors, not only donors and actuators can be used as inputs, but also various kinds of defects, distortions of the crystal circuit at the semiconductor surface, and redundancy of one of the components in a complex semiconductor material. The energy levels in the semiconductor no-go zone, the centers play an important role in the luminescence process. In such cases, luminescence occurs mainly through the centers of recombination.

Electron-hive recombination centers σ_n are characterized by efficient electron and hive σ_p capture efficient cross-sections. Traps and donor inputs have a high probability of pulling electrons out of the conduction zone, but it is almost

impossible for electrons to pass through these centers into the stress zone, because $\sigma_n \gg \sigma_p$. If there is no electron in the center, it is a trap for the electron, and if the center has an electron, the center is the donor.

Transitions during the neutralization of ionized donors and acceptors (Fig. 12.10) can be irradiant in principle and cause radiation in the far infrared region of the spectrum. However, calculations show that the probability of passing by output of phonons is significantly greater than the probability of a radiant transition. For example, for n-type germanium, the capture cross-section for transitions with phonon emission at 4 K is approximately 10^{-12} cm^2 , and for transitions with photon emission the capture cross-section is approximately $4 \cdot 10^{-19} \text{ cm}^2$.

The process of recombination without radiation occurs as follows. First of all, the electron is captured by the excited surface of the donor from the conduction zone. It then makes cascading transitions to the following energy levels, emitting a background in each transition.

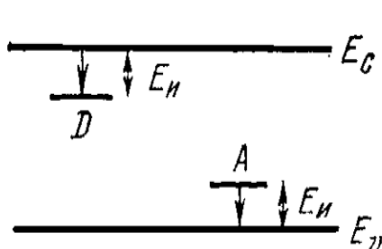


Fig. 12.10. Small jumps

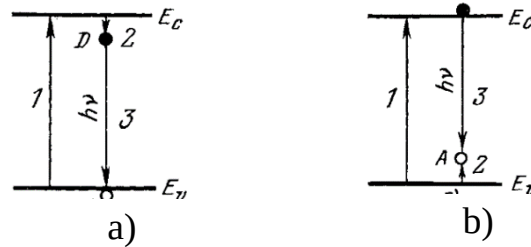


Fig. 12.11. Radiant transitions between zone and entry states

12.4.1. "Zone - access level" transitions

In N- and p-type semiconductors, there may be radiated transitions from the donor state to the stress band, or from the conductivity band to the acceptor layer, as shown in Fig. 12.11. Such centers differ from traps only in that they have a slightly greater depth of layers and are less likely to capture an electron from the permeability zone ($\sigma_n \gg \sigma_p$), although it is of considerable value to catch an electron from the stress zone. According to Fig. 12.11, and after the absorption process (Pass 1), first the electron is captured in the conductivity zone (Pass 2), and then the radiation recombination of the cavity in the localized electron and stress zone occurs (Pass 3).

In another model of luminescence (Fig. 12.11- b), the localized state is located slightly above the stress zone and has a high probability of catching the electron ($\sigma_p \gg \sigma_n$), but there is also a significant probability of simultaneous capture of electrons from the zone of conductivity. Luminescent radiation occurs when a free electron in

the conduction zone is recombined with the swell trapped by this shaft (Transition 3).

In the right transitions, a photon with the following energy is released:

$$h\nu = E_g - E_i \quad (12.7)$$

In case of false transitions, a photon with the following energy is released:

$$h\nu = E_g - E_i - E_p. \quad (12.8)$$

where E_i is the ionization energy of the input atom; E_p is the energy of the phonon involved in recombination.

This type of irradiation recombination is possible only for semiconductors with small concentrations. An increase in the concentration of donors (or actors) causes the energy level of the input to deteriorate and become an access zone that merges with the edges of the main zone.

As an example, Fig. 12.12 shows the cathodoluminescence emission spectrum of zinc-doped gallium arsenide. As can be seen from the figure it is clear that with an increase in the concentration of the receiver (indicated by figures in the curves), an enlargement of the radiation strips occurs, which is caused by narrowing of the forbidden zone due to a high degree of leprocity.

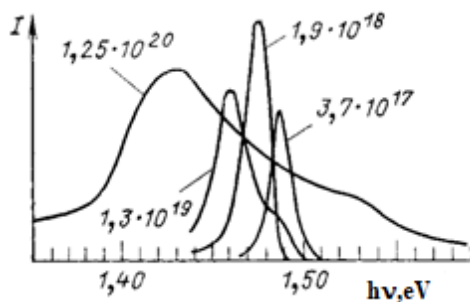


Fig. 12.12. Spectrum of cathodoluminescence emission of zinc-doped gallium arsenide at 4.2 K (cm^{-3})

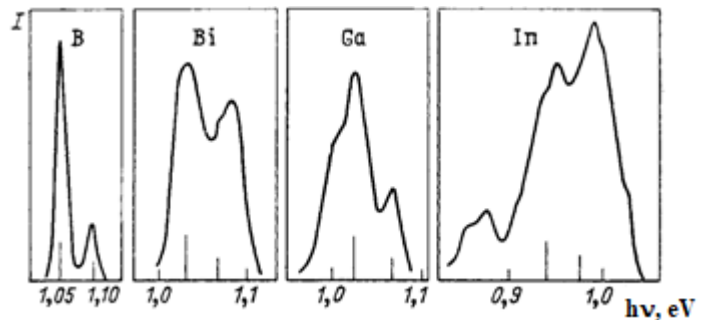


Fig. 12.13. Spectral distribution of input radiation recombination in silicon at 20K

In semiconductors with a non-conductive bandwidth, a phonon must be present in "donor-stress zone" transitions as well as "conduction zone-receptor" transitions in a small concentration of donors. Fig. 12.13 shows the radiation spectra for silicon samples alloyed with different inputs.

A transverse acoustic phonon TA with an energy of 0.016 eV and a transverse optical phonon TO with an energy of 0.055 eV participate in recombination transitions that ensure the preservation of the quasi-impulse of charge carriers. In the picture, a long vertical line $h\nu = E_g - E_i - E_p$ corresponds to the energy

(TO), short $h\nu = E_g - E_i - E_p$ for (TA) energy. Ionization energy values (in electron volts) for inputs are: 0.046 (B); 0.069 (Bi); 0.071 (Ga) and 0.16 (In). In the sample where indium is introduced, phonon-free transitions are observed in the low-energy part of the spectrum.

12.4.2. Donor-akseptor juftlari

If a semiconductor has donor and acceptor inputs that have ionization energy E_d and E_a , respectively, they can form pairs that play the role of motionless molecules in the crystal. The Coulomb interaction between the donor and the acceptor located in a medium with a relative dielectric absorption ϵ_r reduces their bonding energy by the following amount:

$$\Delta E = e^2 / \epsilon_r r, \quad (12.9)$$

where r is the distance between the donor and the acceptor that makes up the pair.

Hence the energy range between the donor and acceptor levels

(Fig. 12.14) would be equal to:

$$h\nu = E_g - E_d - E_a + \frac{e^2}{\epsilon_r r}. \quad (12.10)$$

Because the inputs can occupy discrete states in the crystal lattice (e.g., the substitute input occupies the node), the distance in the donor-acceptor pair varies into finite parts. This should lead to the formation of a thin structure on the spectrum. At great distances ($r > 4.0 \text{ nm}$) between the donor and the acceptor, the radiation stripes stick to each other, forming a wide spectrum. The spectral state of a broad-line and narrow-line structure must depend on the type of inputs. In this case, pairs of different types can be formed, depending on the places occupied by the entrances. If the donor and the acceptor occupy the same places on the grid, then type I couples are formed. If the donor and the acceptor occupy different places, then type II couples are born.

Donor-acceptor pair radiation has been detected in several semiconductors.

Specifically, Fig. 12.15 shows the 1.6 K spectrum of gallium phosphide, in which the donor-acceptor pairs are Si-S and Si-Te in the acceptor Si. Both donors and acceptors take the place of phosphorus at the node

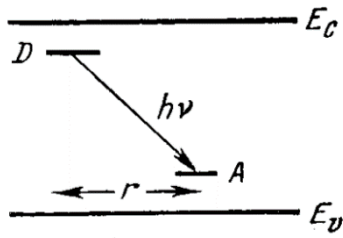


Fig. 12.14. Donor-Acceptor Transitions

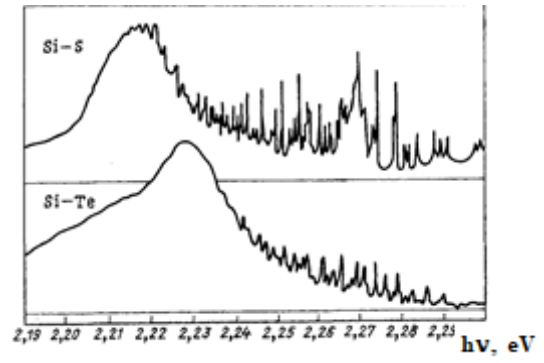


Fig. 12.15. Photoluminescence spectra of GaP crystals at 1.6 K corresponding to donor-receptor recombination in Si-S and Si-Te pairs

12.5. Relaxation of semiconductor luminescence

The semiconductor radiation does not disappear instantly when the excited source is turned off. Since the intensity of the luminescence is proportional to the number of recombination acts per second

$$I_l \sim -\left(\frac{dn}{dt}\right)_r, \quad (12.11)$$

The change in its brightness after removal of the excitation varies depending on the type of luminescence.

In the case of monomolecular radiation, when absorption and the irradiation process occur at the same center, the decrease in excited electron concentration over time is determined by the ratio of:

$$n = n_0 e^{-t/\tau}, \quad (12.12)$$

n_0 is where $t=0$ is the concentration of excited electrons at the moment; The residence time of the τ -electron at the excitation center level.

Hence the monomolecular luminescence intensity after excitation cessation

$$I_l \sim e^{-\frac{t}{\tau}}, \quad (12.13)$$

that is, τ decreases according to the exponential law with immutable time.

For recombination luminescence, the law of extinction is determined by the degree of its excitation. Consider two cases. In weak excitation, the intensity of recombination according to $\Delta n \ll (n_0 + p_0)$, (7.18) when the concentration of overcharge carriers is significantly less than that of Δn equilibrium, i.e.

$$\frac{d\Delta n}{dt} = -\frac{\Delta n}{\tau}. \quad (12.14)$$

In this case, therefore,

$$I_l = I_0 e^{\frac{-t}{\tau}}, \quad (12.15)$$

where Δn_0 is the concentration of overcharge carriers, while I_0 is the luminescence intensity at time $t=0$, i.e. when the excited light is turned off.

A type of semiconductor radiation, expressed in intensity (12.15), is called monomolecular recombination radiation. From the exponential curve of its relaxation, it is possible to determine the time of decomposition of the unbalanced charge carriers τ .

In the case of strong excitation, i.e., $\Delta n \gg (n_0 + p_0)$, the intensity of recombination will be proportional to the square of the concentration of excess charge carriers, i.e.:

$$\frac{d\Delta n}{dt} = -\gamma_r \Delta n^2. \quad (12.16)$$

Whereas the concentration of excess charge carriers decreases over time:

$$\Delta n = \frac{\gamma_r \Delta n_0}{1 + \gamma_r \Delta n_0 t}, \quad (12.17)$$

The intensity of the luminescence is expressed as

$$I_l = \frac{\gamma_r^2 \Delta n_0^2}{(1 + \gamma_r \Delta n_0 t)^2}, \quad (12.18)$$

and bimolecular recombination is called luminescence.

12.6. Influence of temperature on semiconductor luminescence

In semiconductors, the intensity of the radiation decreases with temperature increase, i.e. the luminescence decreases (dwindles) as the temperature increases. If P_r is the probability of radiation-free transitions perceived as being temperature-independent, and P_{nr} is the probability of radiation-free transitions, then the quantum output of luminescence is characterized by η expression.

$$\eta = \frac{P_r}{P_r + P_{nr}}. \quad (12.19)$$

The probability of P_{nr} depends on temperature

$$P_{nr} = P_{nr0} e^{-E^*/kT}. \quad (12.20)$$

In this expression, E^* is the activation energy; P_{nr0} is a temperature-dependent coefficient. Therefore, the temperature dependence of the radiation quantum output has the following form:

$$\eta = \frac{1}{1 + C e^{-E^*/kT}}, \quad (12.21)$$

bu yerda $C = P_{nr0}/P_r$ -konstanta.

The failure of monomolecular luminescence of solids can be explained as follows: when the absorption of light and the radiation occur in the same center. When the centrifugal electron in the absorption of the light quantum is moving to the excited state, at high temperatures the oscillations corresponding to the nearest location of the potential curves U_o and U_b can reach an energy level (point a in Fig. 12.1). Then it switches to the main state as a result of recombination without radiation. This is illustrated by an jump-off bar in Fig. 12.1. Luminescence is characterized by the formula (12.21) for the mechanism of luminescence atuation through temperature. In this case, the activation energy required to transfer the electron from the E^* -excited state to the nearest location point of the potential curves U_o and U_b is the activation energy.

The temperature extinction of semiconductor luminescence, which results from the recombination of free charge carriers through a localized center, is also characterized by (12.21). For example, the center of radiation is located in the near distance to the ceiling of the stress zone (Fig. 12.11, b). Such a center can be excitation by absorbing a quantum of light (see Fig. 11.31, a). The ionized center then captures the electron from the stress zone. An electron snatched away by centrifugal light cannot return to the center by radiation, but it is forced to diffuse for some time in the conductivity zone until it can recombine without radiation with the cavity in the stress zone. For example, non-radiation-free transitions containing germanium copper have a thermal activation energy of $E^*=0.14$ eV.

In a radiative recombination of free electrons and free swells, the dampening of the luminescence under the influence of temperature

$$\eta = \frac{1}{1+CT^g} \quad (12.22)$$

It can be described by the formula. Here C and g are constants.

To obtain this bond, the theory of recombination of an electron in the conductivity zone and a pit in the stress zone without correct radiation was used.

12.7. Spontaneous (self-absorbed) and forced radiation from the atom

Before proceeding to the description of the conditions under which laser radiation occurs in semiconductors, let us consider the process of absorption and emission of light in an atom. Simpler than the whole set of quantum states of an atom, consider two states with only E_1 and E_2 energy (Fig. 12.16).

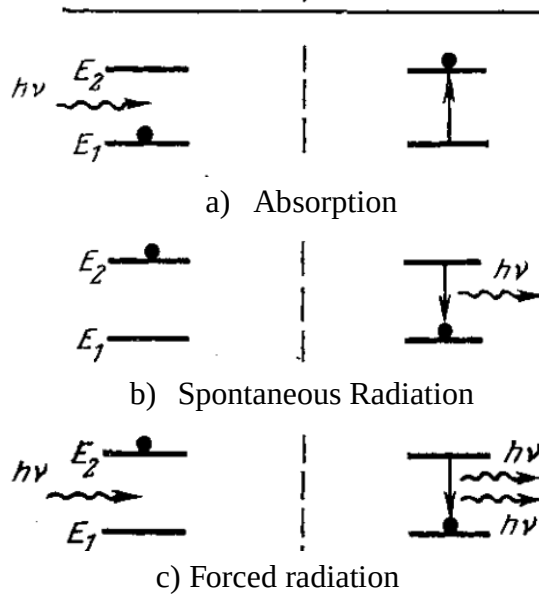


Fig. 12.16. There are three main types of transitions between energy levels E_1 and E_2 : black dots indicate the state of an atom. The initial case corresponds to the picture on the left; The final state after the relevant process has occurred is shown on the right side

Let the spectral energy density of such an atom be affected by the radiation in which it u_ν is. Anv unexcited atom that absorbs the radiation energy h passes into an excited state (Fig. 12.16, a). Moreover, the larger it u_ν is, the more often transitions with absorption occur. If N_1 is the number of atoms in the unexcitation state, and it $u_\nu B_{12}$ is the probability of passing by absorption, then the number of transitions by absorption in units of time in unit volume is equal to:

$$dN_{12} = B_{12}u_\nu N_1 \quad (12.23)$$

As a result of the absorption of the quantum of light, the atom, which is in an excited state, returns to the normal state from the excited state after a certain time, and $E_2 - E_1$ releases excess energy in the form of radiation $h\nu$. A radiative transition can occur spontaneously, without dependence on the influence of the external field (Fig. 12.16, b). This transition is called spontaneous, and the radiation is called spontaneous (spontaneous) radiation.

Spontaneous jumps are random, and their probability is determined by A_{21} . If the number of atoms in the N_2 -excited state is the number of spontaneous transitions per unit of time, then the number of spontaneous transitions per unit of time per unit of volume is:

$$dN'_{21} = A_{21}N_2. \quad (12.24)$$

Therefore, the decrease in the number of excited atoms as a result of spontaneous transitions in dt time is as follows:

$$-dN_2 = A_{21}N_2 dt. \quad (12.25)$$

Integrating this equation, we get the law of decreasing over time of the number of atoms excited during the spontaneous transition:

$$N_2 = N_{20}e^{-A_{21}t}, \quad (12.26)$$

where N_{20} is the number of atoms in the excited state at time $t = 0$.

Since energy $E_2 - E_1 = h\nu$ is released in each spontaneous transition, dN'_{21} the following amount of energy is released per unit of time as a result of the transitions:

$$h\nu dN'_{21} = h\nu A_{21}N_2. \quad (12.27)$$

The intensity of spontaneous radiation (energy emitted in units of time) is equal to:

$$I = h\nu A_{21}N_{20}e^{-A_{21}t} = I_0e^{-A_{21}t}, \quad (12.28)$$

here

$$I_0 = h\nu A_{21}N_{20}. \quad (12.29)$$

$A_{21}N_2dt$ Since this is the number of atoms that have completed the transition to 2-1 in dt time, t determines the number of atoms that have "lived" in the perturbed state over time. Thus, ∞ the duration of remaining in the excited state of all atoms that have experienced a spontaneous transition for a period of time from 0 to ∞ will be:

$$T = \int_0^{\infty} t A_{21} N_2 dt = \int_0^{\infty} N_{20} A_{21} t e^{-A_{21}t} dt = \frac{N_{20}}{A_{21}}. \quad (12.30)$$

From this it follows that the average residence time of an atom in the excited state

$$\tau = \frac{T}{N_{20}} = \frac{1}{A_{21}} \quad (12.31)$$

So so.

In accordance with this (12.28), the formula can be written as follows:

$$I = I_0 e^{-t/\tau}. \quad (12.32)$$

Thus, once excitation is stopped, the intensity of spontaneous radiation decreases exponentially with constant τ time over time.

Since the spontaneous transitions are random, the spontaneous radiation is of a statistical nature, it is non-coherent, that is, photons are generated, which, although fixed in frequency, scatter in voluntary directions and have voluntary phases. This is because, by their random nature, spontaneous acts of radiation occur independently of each other at different moments in time.

However, an atom in the excited state (E_2 state) can return to its normal state (E_1 state) under the influence of the radiation field. Such a transition is called compulsory, and radiation is called forced, induced, or stimulated (Fig. 12.16-v).

Induced jumps outweigh the energy density of the excited light, the u_ν faster it occurs. Therefore, the probability of these transitions for a unit of time is that it is $u_\nu B_{21}$. If N_2 is the number of atoms in the excited state, then the number of induced transitions for a unit unit of time is:

$$dN_{21}'' = u_\nu B_{21} N_2. \quad (12.33)$$

Consequently, ν the frequency of external radiation results not only in absorption of the light, but also in additional, induced radiation.

Thus, the average number of gradient transitions per unit of time is determined by spontaneous and forced transitions and is equal to:

$$dN_{21} = dN_{21}' + dN_{21}'' = (A_{21} + u_\nu B_{21}) N_2. \quad (12.34)$$

In the stationary state, based on the detailed equilibrium principle, the number of correct and inverse processes in a unit of time will be the same. So,

$$dN_{12} = dN_{21} \quad (12.35)$$

or

$$(A_{21} + u_\nu B_{21}) N_2 = u_\nu B_{12} N_1. \quad (12.36)$$

From this we find:

$$\frac{N_2}{N_1} = \frac{u_\nu B_{12}}{A_{21} + u_\nu B_{21}}. \quad (12.37)$$

Since equilibrium has a statistical property, the ratio N_2/N_1 can be expressed using the Boltzman distribution, according to which the energy is $E_1, E_2, \dots, E_m, \dots$. For a system that has states that are, the probability that it is in the energy state E_m is:

$$W_m = C g_m e^{-E_m/kT}, \quad (12.38)$$

Here g - is the statistical weight that takes into account the state m . So, we can write:

$$\frac{N_2}{N_1} = \frac{g_2 e^{-E_2/kT}}{g_1 e^{-E_1/kT}} = \frac{g_2}{g_1} e^{-(E_2-E_1)/kT}. \quad (12.39)$$

A comparison of (12.37) with (12.39) follows the following conclusion:

$$\frac{u_\nu B_{12}}{A_{21} + u_\nu B_{21}} = \frac{g_2}{g_1} e^{-(E_2-E_1)/kT}, \quad (12.40)$$

From here we get, taking into account $E_2 - E_1 = h\nu$:

$$u_\nu = \frac{g_2 A_{21}}{g_1 B_{12} e^{h\nu/kT} - g_2 B_{21}}. \quad (12.41)$$

We use the following boundary conditions to determine the coefficients A and B . When the $\rightarrow \infty$ energy density is t , it u_ν must tend to infinity, i.e., the intensity in the expression (12.41) tends to zero. From this follows:

$$g_1 B_{12} = g_2 B_{21}. \quad (12.42)$$

Taking this into account (12.41), the formula can be written as follows:

$$u_\nu = \frac{g_2 A_{21}}{g_1 B_{12}} \frac{1}{e^{h\nu/kT} - 1}. \quad (12.43)$$

At the first approximation, we get the following -- $e^{h\nu/kT} = 1 + h\nu/kT$. From this

$$u_\nu = \frac{g_2 A_{21} kT}{g_1 B_{12} h\nu}. \quad (12.44)$$

The Raleigh-Jeans formula applies:

$$u_\nu = \frac{8\pi\nu^2}{c^3} kT. \quad (12.45)$$

Comparing (12.44) and (12.45), we get:

$$\frac{g_2 A_{21}}{g_1 B_{12}} = \frac{8\pi h \nu^3}{c^3}. \quad (12.46)$$

See with others:

$$A_{21} = \frac{8\pi h \nu^3}{c^3} \frac{g_1}{g_2} B_{12} = \frac{8\pi h \nu^3}{c^3} B_{21}. \quad (12.47)$$

Thus, in the forced transition, the direction and phase inducing the recombination are generated by the same particle, the photon, as the photon.

Now let's look at the conditions under which induced radiation can be observed. Let there be a parallel beam of monochromatic light with intensity I . When passing through a layer of material of thickness Δx , the intensity of light decreases by the amount of $-\Delta I = \alpha I \Delta x$. This decrease in luminous intensity is determined by the fact that in atoms it occurs with the absorption of energy in the amount $u_\nu B_{12} h \nu N_1 \Delta x$ and it occurs induction transitions with energy radiation equal to $u_\nu B_{21} h \nu N_2 \Delta x$. As a result, we can write:

$$-\Delta I = \alpha I \Delta x = (B_{12} N_1 - B_{21} N_2) u_\nu h \nu \Delta x. \quad (12.48)$$

As can be seen from (12.48), absorption and induced radiation are equal to the formula, but entered by different signs. If the first dose indicates normal absorption, it is considered a "positive" absorption, then induced radiation can be considered a "negative" absorption.

If we consider that a parallel beam of light with a cross-section of 1 cm^2 propagates at the speed c , then the intensity of the radiation is equal to the following:

$$u_\nu = \frac{1}{c}, \quad (12.49)$$

Then the coefficient of absorption according to the equations (12.42), (12.48) and (12.49) α takes the following form:

$$\alpha = \frac{h\nu}{c} (B_{12} N_1 - B_{21} N_2) = \frac{h\nu}{c} \left(1 - \frac{g_1}{g_2} \frac{N_2}{N_1}\right) B_{12} N_1. \quad (12.50)$$

Considering (12.39):

$$\alpha = B_{12}N_1(1 - e^{-(E_2-E_1)/kT}) \frac{h\nu}{c}. \quad (12.51)$$

Since $E_2 > E_1$ is, the exponential term in (12.51) is always less than one, and the magnitude in parentheses is positive. This means that at any final temperature the coefficient of absorption of the aggregate at the thermodynamic equilibrium condition is always positive, i.e. $\alpha > 0$. In this case, the usual positive absorption of light occurs, and induction jumps only reduce it. This is because in the equilibrium state, the number of unexcited atoms will always be greater than the number of atoms excited, i.e., $N_1 > N_2$.

From the Boltzman distribution (12, 39) it is evident:

$$T = \frac{E_2 - E_1}{k \ln \frac{g_1 N_2}{g_2 N_1}}. \quad (12.52)$$

For example, if $E_2 > E_1$ is equilibrium and $g_1 N_2 < g_2 N_1$, then $T > 0$. However, the following

$$g_1 N_2 > g_2 N_1 \text{ or } \frac{N_2}{N_1} > \frac{g_2}{g_1}, \quad (12.53)$$

If the inequality is met, it would be < 0 based on (12.50). α Thus, in order to create a medium with negative absorption, it is necessary to achieve an equilibrium state in which the number of excited atoms is greater than the number of atoms in the normal, unexcited state. (12.53) The busyness of the energy levels of the atoms satisfying the inequality is called inverse busyness.

12.8. Forced radiation of solids

Now let's look at specific ways to generate inverse employment. It is important to note that in a two-tier system, it is impossible to generate a stable inverse bandage. If particles are emitted by resonant radiation to a really high level, then $N_2 g_1 = N_1 g_2$, and the probabilities of the radiation processes are equalized, and inverse employment becomes impossible. This can happen when particles to the high energy level are emitted by inter-collision (e.g. by electrons or atoms). In this case, as the number of particles in the upper layer increases, the probability of particles falling into the lower layer as a result of collision with other particles also increases and does not exceed the amount defined by the Boltzman equation ($N_2 = N_0 \exp(-\varepsilon_2/kT)$). Since the temperature of these particles is always positive, the condition of $N_2 < N_1$ is fulfilled and the inverse occupation is not formed. In order for inverse occupancy to be formed in this case, the processes for ejection and fall of particles to the high ground must be diverse.

For inverse employment, there must be at least three energy lines in this system. One of these should be the main line. The simplest method of generating inverse employment, i.e., the three-tier system method, is shown in Fig. 17. Inverse employment is formed between the metastable 2 lines and the main layers. The particle saturation of the metastable 2 layer is formed as a result of the transition of particles from the 3 layers by non-radiation-free relaxation with a probability of S_{32} to the 2nd layer. To simplify the process being seen, let's ignore other non-radiant transitions. This does not affect the results obtained. If we define the correct and inverse transitions and probabilities between lines 1 and 3 and 1 and 2, respectively, W_{13}, S_{31} W_{12} va W_{21} , then in the equilibrium state we can determine the balance of particles in these layers as follows

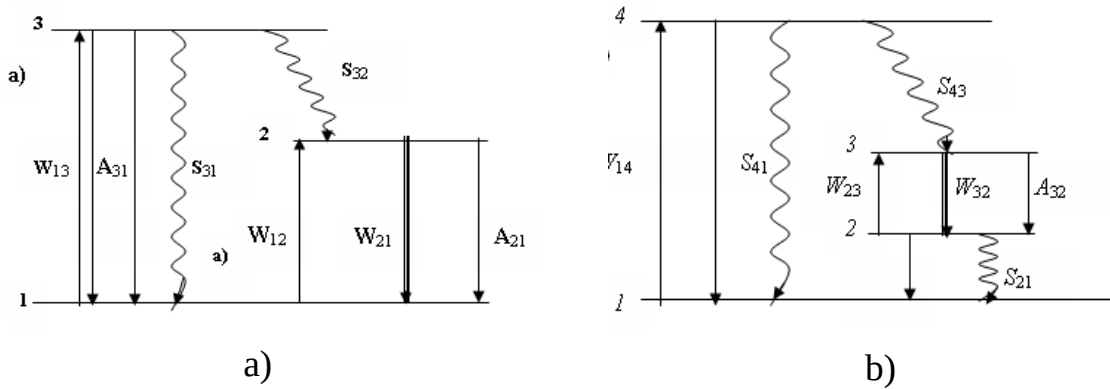


Fig. 12.17. An inverse employment generation scheme in a three(a) and four(b) quantum system.

$$\left. \begin{aligned} \frac{dN_3}{dt} &= W_{13}N_1 - (W_{31} + A_{31} + S_{32})N_3 = 0 \\ \frac{dN_2}{dt} &= W_{12}N_1 - (A_{21} + W_{21})N_2 + S_{32}N_3 = 0 \end{aligned} \right\} \quad (12.54)$$

va

$$N_0 = N_1 + N_2 + N_3 \quad (12.55)$$

It is possible to write in views. Here N_0 —is the full concentration of active particles. N_1, N_2, N_3 are the – concentrations of particles at the corresponding states, A_{31} and A_{21} —are the probabilities of spontaneous transitions. Equations (12.54) and (12.55) above show that the relative saturation of levels 1 and 2 is as follows

$$\frac{N_2}{N_1} = \frac{W_{13}(S_{32}/(W_{31} + A_{31} + S_{32})) + W_{12}}{A_{21} + W_{21}} \quad (12.56)$$

It can be identified by expression. The effective filling of the upper 2nd floor is as follows:

$$A_{31} \ll S_{32}, \quad W_{31} \ll S_{32} \quad (12.57)$$

This can only happen if the conditions are met.

In this case (12.56) the expression follows

$$\frac{N_2}{N_1} \approx \frac{W_{13} + W_{12}}{A_{21} + W_{21}} \quad (12.58)$$

It takes the view. If we subtract 1 from both sides of this parity, then

$$\frac{N_2 - N_1}{N_1} \approx \frac{W_{13} - A_{21}}{A_{21} + W_{21}} \quad (12.59)$$

When the above conditions are met $N_3 \rightarrow 0$ va $N_1 + N_2 \approx N_0$ So. If levels 1 and 2 are not separated ($W_{12} = W_{21}$), 1 layer infill

$$N_1 = \frac{N_0(A_{21} + W_{21})}{A_{21} + W_{13} + 2W_{12}} \quad (12.60)$$

It will be equal to the magnitude, if we put N_1 (12.59) into the expression and make the modifications

$$N_2 - N_1 = N_0 = \frac{W_{13} - A_{21}}{A_{21} + W_{13} + 2W_{12}} \quad (12.61)$$

Ifodani olamiz.

As can be seen from this formula (12.59), the inversion is as follows only if the conditions (12.59) are met, i.e. the probability of 3 layers being filled is greater than the probability of spontaneous radiation transition of particles from 2

$$W_{13} > A_{21} \quad (12.62)$$

Without it. In the three-dimensional system of consideration, it is possible to introduce the chromium ion in a ruby (ruby) crystal as the active particle in a laser.

The main disadvantage of inverting inverse employment in a three-layer system is the condition that a large number of particles are agitated. Really $N = N_1 + N_2$ bo'lib, invers bandligi $N_2 > N_1$ The following conditions shall be fulfilled $N_0 - N_1 > N_1$ On the look, if we write it down, then

$$N_2 > \frac{N_0}{2} \quad (12.63)$$

It originates.

Thus, in order to obtain inverse occupancy in a three-layer system, more than half of all particles must be located in the upper layer. This condition places significant demands on the rate of excitation, as well as on the throttle power.

One of the most common methods free of these drawbacks is the four-layer system, the diagram of which is shown in Fig. 17b. In such an energy system, increased radiation occurs when active particles give compulsory radiation from 3 to 2 kt. The particle filling of the 3rd layer occurs as a result of the transition of particles from the 4 th layer to the 3 rd layer without radiation, as shown in Figure 17B. The rapid release of the 2 and layer from the particles occurs by the transition

of the particles in it to the 1st layer without radiation. The condition for forming a stable inverted employment situation in such a system are:

$$W_{23} > A_{32}, W_{32} < S_{43}, S_{21} > W_{32} \quad (12.64)$$

Thus, since the lower working line (i.e., the 2nd layer) is not the main, the inverse occupancy will not depend on the relative filling of the 2 and 3rd floors with the particles, just as in the three-line system, the upper working line does not have to be filled with a large absolute value. In this conditional environment, inverse bandage allows for even at smaller throttle capacities. Examples of this method include obtaining inverse bandages in CO_2 , CO molecules in gas lasers, neodymium ion in solid-state lasers, etc.

It turns out that the methods of inverse formation depend not only on the system of concrete energy layers, on the properties of the particles, on the other components that make up the active medium. As active mediums of lasers, they can be mixtures of gases, various condensed substances, crystals, glasses, semiconductors and liquids. In laser systems, optical, gaseous, gasodynamic and chemical tincture methods are used.

Control Questions

1. List and explain the types of luminescence
2. Describe the monomolecular radiation of solids
3. How does recombinant radiation of semiconductors occur in fundamental transitions?
4. Describe a Highly Alloyed Semiconductor
5. Explain excitone recombination
6. How does recombination radiation occur during transitions between zones and entry layers?
7. How are donor-receptor pairs formed?
8. What is the relaxation of semiconductor luminescence?
9. How does temperature affect the luminescence of semiconductors?
10. Substantiate the differences between spontaneous and forced radiation of the atom
11. Explain the forced radiation of solids

Chapter XIII .

PHOTOELECTRIC PHENOMENA IN SEMICONDUCTORS

13.1. Internal photo effect and photoconductivity

The phenomenon of the intrinsic photoeffect in semiconductors can be explained as follows. When a light beam hits a semiconductor material, its electrical resistance changes, such a phenomenon is called the photoresistive effect. The electrical extra conductivity caused by the charge carriers generated by optical generation is called photoconductivity.

In order for the internal photoeffect phenomenon to be observed, it is necessary first to absorb the light in a crystal. Depending on the nature of light absorption, it is possible to distinguish between private or ingressal photoconductivity. In private photoconductivity, interband absorption of light is observed. In this case, the absorbed photon energy for a straight-pass semiconductor should not be less than the bandwidth of the bandwidth $h\nu \geq E_g$ (optical transition 1 in Fig. 13.1):

$$h\nu \geq E_g \quad (13.1)$$

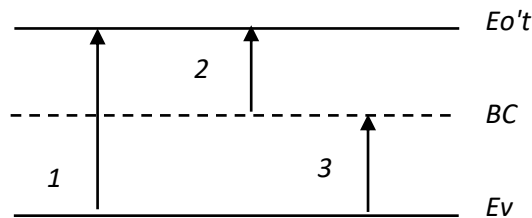


Fig. 13.1. A scheme of optical transitions that generate photoconductivity.

When the law of conservation of a quasi-impulse in a non-equilibrium transition semiconductor applies, the energy of the longitudinal (red) boundary of photoconductivity is E_p :

$$h\nu = E_g + E_p \quad (13.2)$$

If there are local energy levels in the exclusion zone, the optical transitions may be between the main zones and the access layers (optical transitions 2 and 3 in Fig. 13.1. The photoconductivity in this state is called inlet photoconductivity.

The phenomenon of exciton absorption of a photon is also observed in semiconductor materials. In such absorption, the generated electron and the pit are in a state connected to each other, that is, in a neutral state, and does not initially form a photocurrent. However, excitons can form a photocurrent by interacting with different entry centers, defects, and phonons during the movement in the crystal.

Photon energy can also be absorbed by free charge carriers, but the concentration of charge carriers does not increase due to this. If the concentration of charge carriers increases strongly, the kinetic energy of the charge carriers will increase in the process of absorption of light, and at the expense of such energy, new charge carriers may be awakened.

To understand photoconductivity, we take the following approach. In cases where the (excessive) non-equilibrium electrons Δn and the spikes Δp have more energy than the energy of the balanced charge carriers Δp , due to optical generation, such charge carriers also have the same energy and quasi-impulse distribution as the main charge carriers. As a result, the conductivity of the semiconductor is determined by the contributions of the main charge carriers n_o and p_o and the photogenerated charge carriers Δn and Δp :

$$\sigma = q (n_o + \Delta n) \mu_n + (p_o + \Delta p) \mu_p \quad (13.3)$$

For the conductivity when the light does not fall is as follows:

$$\sigma = q (n_o \mu_n + p_o \mu_o) \quad (13.4)$$

And for the conductivity when the light falls, the following will be appropriate:

$$\sigma_F = \sigma - \sigma_o = q (\Delta n \mu_n + \Delta p \mu_p) \quad (13.5)$$

The concentration of photogenerated charge carriers Δn and Δp depends on the incidence light intensity and wavelength. In this case, for the absorbed light energy in units of time and units of volume, it is possible to write the following:

$$dI/dx = \alpha I \quad (13.6)$$

For the number of photogenerated particles in units of time and units of volume, we can write:

$$G = \beta \alpha I / h\nu \quad (13.7)$$

Here β is the photoionization quantum output indicator.

When light of the same intensity hits a semiconductor, the recombination rate also increases, equaling the rate of generation. In this case, a stable photoconductivity is formed, for which the following formula is appropriate:

$$\sigma_{F,t} = \beta \alpha q (\mu_n \tau_n + \mu_p \tau_p) / h\nu \quad (13.8)$$

For the specific photosensitivity of a semiconductor, the following expression would be appropriate:

$$SF = \sigma F / I. \quad (13.9)$$

If the value of the 1st or 2nd in the formula (13.5) is very large relative to the second, i.e., the contribution of charge carriers of one type is large, then the photoconductivity in this case is called monopolar photoconductivity. The photoconductivity formula for such a case is as follows:

$$\sigma_{F,t} = \beta \alpha q \tau \mu / h\nu \quad (13.10)$$

The following formula is appropriate for stationary photocurrent density:

$$J_F = \sigma_{F,t} \mathcal{E} = [\beta \alpha q (\mu_n \tau_n + \mu_p \tau_p)] \mathcal{E} / h\nu \quad (13.11)$$

If we take into account that $\mu_n \mathcal{E} = v_n$ and $\mu_p \mathcal{E} = v_p$, and assuming that the semiconductor size in the field direction is l , and also taking into account $t_n = l / v_n$ and $t_p = l / v_p$ (13.11), the equation can be written as follows:

$$- \Delta p / \tau_p + D d^2 \Delta p / dx^2 + G e^{-\alpha x} = 0 \quad (13.13)$$

By paying attention to the light intensity of constant photoconductivity, 2 different states can usually be seen. At low radiation intensities, a straight line coupling is observed and at large irradiation, sublinear bonding.

13.2. Photoconductivity in the presence of surface recombination and diffusion of charge carriers

In the previous paragraph, the state of homogeneous photogeneration across the semiconductor volume was seen in the interpretation of photoconductivity, and surface recombination was not taken into account.

Now let's see a case where the light falling on the surface of a semiconductor is unevenly absorbed in thickness and there is a diffusion of charge carriers.

In order to calculate the constant concentration of unbalanced charge carriers in the absence of retaining (adhesion) centers, it is necessary, for example, to solve the following continuity equation for holes:

$$- Dp/t_p + D d^2 \Delta p / dx^2 + G e^{-\alpha x} = 0 \quad (13.13)$$

In this equation, 1-had represents recombination, 2-had represents diffusion, and 3-had represents generation:

$$G e^{-\alpha x} = [\beta (l - R_s)^2 I_0 / h\nu] e^{-\alpha x} \quad (13.14)$$

where β is the quantum output, R_s is the return coefficient of the semiconductor surface.

If we assume that light of I_0 intensity is falling on the surface of a semiconductor, the boundary condition for the surface $x=0$ can be written as:

$$(Dd\Delta p/dx)_{(x=0)} = s\Delta p_{(x=0)} \quad (13.15)$$

where s is the rate of recombination on the illuminated surface. A condition is not required for the back surface of a semiconductor of thickness d , since $d \gg L$ and $d \gg 1/\alpha$.

(13.13) The general solution of the equation is as follows:

$$\Delta p(x) = Ae^{-x/L} - (\tau Ge^{-\alpha x} / L^2 \alpha^2 - 1) \quad (13.16)$$

If we take into account the marginal conditions, we get:

$$\Delta p(x) = (G\tau / \alpha^2 L^2) [(\alpha L^2 + s\tau) / (L + s\tau)] e^{-x/L} - e^{-\alpha x} = 0 \quad (13.17)$$

If we integrate the latter equation from 0 to d , we can determine the number of fully generated holes.

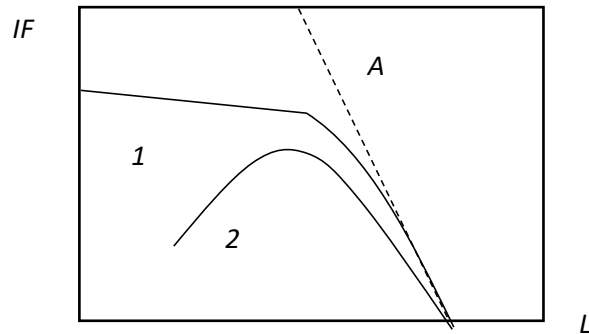


Fig. 13.2. Dependence of the photocurrent on the light wavelength for the presence and absence of surface recombination

Looking at the spectral characteristic presented in Fig. 13.2, it can be understood that in the absence of recombination, the photoc value decreases and saturates with increasing wavelength ($s = 0$; curve 1). The spectral line has a maximum value point (curve 2) when a recombination process is present.

13.3. Dember effect. Photoelectromagnetics effect

When light falls on the surface of a semiconductor corresponding to the private absorption range, light absorption and the formation of electron-cavity pairs are observed near the surface. Such photogenerated particles are diffused towards

the intense. Since the behavior of the electrons is higher than in the swells, they are diffused further (deeper) (Fig. 13.3).

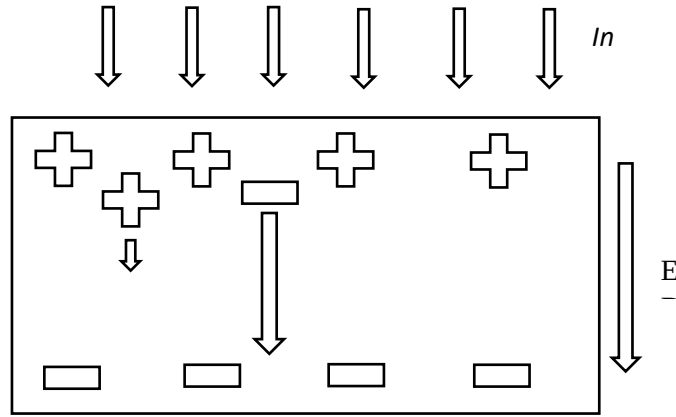


Fig. 13.3. Appearance of the Dember electric driving force in a semiconductor

As a result, a positively charged state is manifested on the surface relative to the volume. The resulting charge distribution is referred to σ surface charge density.

Assuming that there are no surface and volume recombination centers in a semiconductor, we define the Dember EYK. In this case, it is necessary to solve the following equations for the electron and the holes:

$$J_n = q (n_o + \Delta n) \mu_n \mathcal{E}_D + q D_n d\Delta n/dx \quad (13.18)$$

$$J_p = q (p_o + \Delta p) \mu_p \mathcal{E}_D + q D_p d\Delta p/dx \quad (13.19)$$

For the steady state the current is equal to 0:

$$J = J_n + J_p = q [(n_o + \Delta n) \mu_n + (p_o + \Delta p) \mu_p] \mathcal{E}_D + q (D_n dn/dx - D_p dp/dx) = 0 \quad (13.20)$$

We also get the following for Dember EYK:

$$\mathcal{E}_D = (-kT/q) [(\mu_n - \mu_p) d\Delta n/dx] / (n\mu_n - p\mu_p) \quad (13.21)$$

where $n = n_o + \Delta n$ and $p = p_o + \Delta p$. If we take into account $b = \mu_o/\mu_p$, we get:

$$\mathcal{E}_D = (-kT/q) [(b-1)/(b+1)] \ln(1/\sigma) (d\Delta\sigma/dx) \quad (13.22)$$

$$U_D = (-kT/q) [(b-1)/(b+1)] \ln[1+(b+1)\Delta n/(bn_o+p_o)] \quad (13.23)$$

The value of the Dember EIC depends on the difference in the behavior of the electrons and the holes. If they are mutually equal, the Dember voltage is equal to 0.

About the photoelectromagnetic effect, we need to know the following. If a semiconductor is illuminated with a strong light, and it is placed transversely into

the charge carrier diffusion, an IED occurs, and this phenomenon is called the photoelectromagnetic effect (FEM).

Charge carriers generated by light are affected by the Lorentz force of the magnetic field. This force deviates from the direction of motion of the chargers and the direction of the magnetic field. Circuits and cavities deviate in different directions, and as a result positive and negative charges accumulate in opposite directions, i.e., an electric field or *Core* voltage occurs U_y .

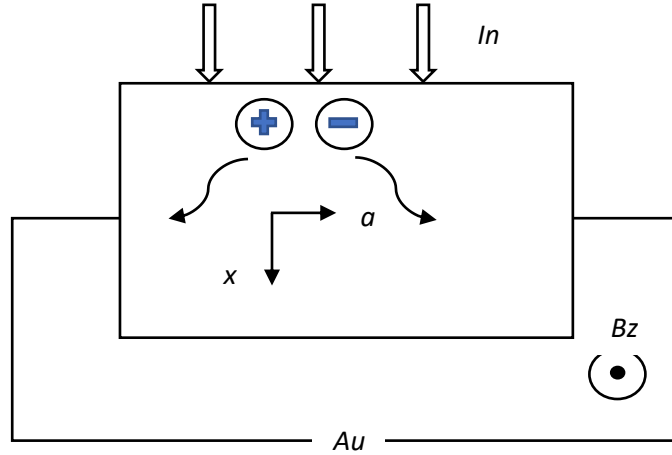


Fig. 13.4. Formation of the IED during the photoelectromagnetic effect

U_y determine the domestic voltage, let a light fall on a semiconductor where surface recombination is very rare, and electron-cavity pairs appear near its surface. For the generated field potential for the weak magnetic field ($\mu B \ll l$), the potential is expressed as:

$$\varphi_n = \mu_n^{\{H\}} B \quad (13.24)$$

$$\varphi_p = \mu_p^{\{H\}} B \quad (13.25)$$

where $\mu_n^{\{H\}}$ and $\mu_p^{\{H\}}$ are behavioral for electrons and swells, respectively.

The following common expressions can be written for a complete diffusion current:

$$J_y = -q r D B (\mu_n + \mu_p) d\Delta n/dx \quad (13.26)$$

where r is the Hall factor $\mu_n = r \mu_n^{\{H\}}$ va $\mu_p = r \mu_p^{\{H\}}$

If we take the size of a semiconductor on the z -axis as b , the following formula for the short-circuit current can be obtained:

$$J_{qt} = -q r D B (\mu_n + \mu_p) b [\Delta n(0) - \Delta n(t)] \quad (13.27)$$

That means that the current density flowing along the surface zx is equal to 0, and the short circuit current flowing in the y direction is determined by the FEM-effect voltage.

$$U_{FEM} = q r a B (\mu_n + \mu_p) \beta L I / h\nu\sigma_o \quad (13.28)$$

where l is the length of the sample, and it is much larger than the free-running path of the charge carriers; A is the size of a semiconductor along the y-axis.

The value of the FEM effect voltage depends on the magnetic field inductance, proportional, and the concentration of charge carriers on the opposite side.

It turns out that the FEM effect voltage depends on the intensity of illumination under low lighting conditions. And with increasing lighting, it tends to saturate.

13.4. P-n Transition Photo Effect. Scotty Barrier Photo Effect

It is known that in the p-n-transition there is a potential barrier due to the diffusion of the main charge carriers due to the electric field. Thermodynamic equilibrium (a) and covered cases (b va c) An energy scheme of the p-n-transition for) is shown in Fig. 13.5.

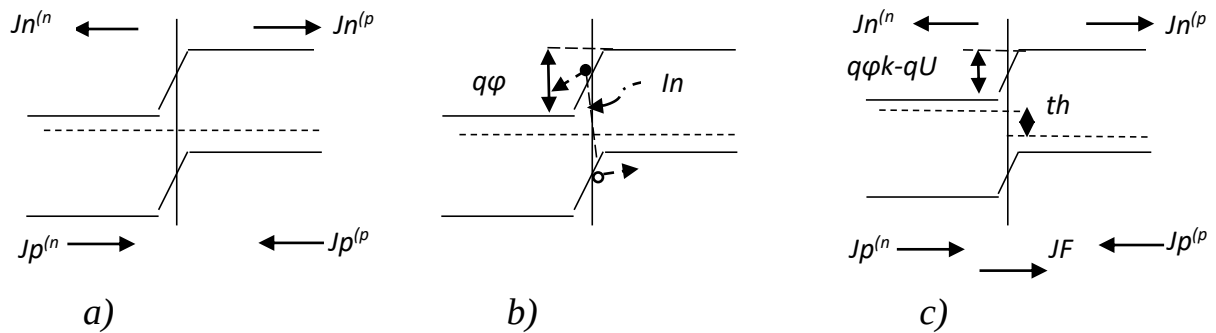


Fig. 13.5. Energy scheme of p-n-transition for the thermodynamic equilibrium (a) and illuminated states (b and c)

For the state of equilibrium of the currents in the outline, it can be written as follows:

$$-J_n^{(n)} + J_n^{(p)} - J_p^{(p)} + J_p^{(n)} = 0 \quad (13.29)$$

For full saturation tobacco, we write:

$$J_s = J_{ns} + J_{ps} = (qD_n n_p / L_n) + (qD_p p_n / L_p) \quad (13.30)$$

When a photon with a greater energy than the bandwidth of the semiconductor is hit (Fig. 13.5 b), an electron-chasm pair is generated. On the permeability zone, the pit appears on the electron, and on the valent zone. Due to the decrease in the potential barrier that occurs when the light does not fall, the electrons are diffused from the p-zone to the n-zone and the spikes from the n-zone to the p-zone. The result is the valve EYK U_{ven} and J_F photocurrent:

$$J = J_F - J_n^{(n)} + J_n^{(p)} - J_p^{(p)} + J_p^{(n)} = 0 \quad (13.31)$$

The density of the nonbasic charge carrier current in the light will be equal to their values at equilibrium state:

$$J_n^{(p)} = J_{sn}^{(p)} ; \quad J_p^{(n)} = J_{sp} \quad (13.32)$$

Due to the decrease of the potential barrier by qU under illuminated conditions, the value of the main charge carrier flux increases as follows:

$$J_n^{(n)} = J_{sn} e^{(qU/kT)} \quad (13.33)$$

$$J_p^{(p)} = J_{sp} e^{(qU/kT)} \quad (13.34)$$

From the above equations we can obtain:

$$U = (kT/q) \ln[1 + (J_F - J) / J_s] \quad (13.35)$$

The resulting equation is called the photodiode equation for any condition. To find the value E_{YK} or the pure operating voltage, we take the state $J = 0$:

$$U_{ven} = (kT/q) \ln[1 + (J_F / J_s)] \quad (13.35)$$

The photo-effect in the Scotty barrier is different from that in the p-n transition. As we know, in the contact of a metal and a semiconductor, a Scotty barrier occurs due to the redistribution of charge carriers on its surface. If the output of electrons from a metal is greater than the output of an n-type semiconductor ($\Phi_M > \Phi_Y$), a Scotty barrier for electrons is formed. If the condition ($\Phi_M < \Phi_Y$) is met, a Scotty barrier for holes will appear. The following formula for the height of the barrier is appropriate:

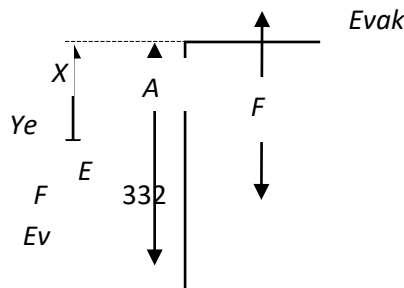
$$q \phi_k = (\Phi_M - \Phi_Y) \quad (13.36)$$

The specificity of the photoeffect observed in the Scotty barrier is that if the photon energy $h\nu > q\phi_k$. In this case, then the electron in the metal passes into the semiconductor by carrying additional energy, overcoming the barrier, where it is negatively charged, photo-E.YU.K is formed. Electrons that are excited from the semiconductor's valent zone to above the Fermi level are tunneled into the metal, but usually their contribution is very small. If $h\nu > E_g$. If the condition is met, the electron-spurs generated around the Scotty barrier are separated in the contact field and photo- E.YU. K forms.

On the basis of the Scotty barrier, usually large-area photocells are prepared. They are used to record photocurrents produced by alpha particles, electrons excited by gamma rays, X-rays, and neutron streams.

13.5. External Photoeffect

Usually, external photoeffect refers to the phenomenon of photoelectron emission on the surfaces of metals. However, as the application of semiconductor materials in micro- and nano-electronics expands, their attention to the external photoeffect is growing. To understand the external photoeffect in semiconductors, we refer to Fig. 13.7.



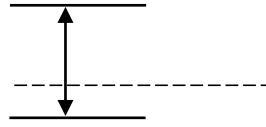


Fig. 13.7. Simplified Diagram of the External Photoeffect in Semiconductors

Using Fig. 13.7, it can be described that the emission of an electron from a semiconductor to a vacuum is called an external photoeffect. E_{vak} is the energy of an electron in which energy is separated from the semiconductor and has a kinetic energy of 0 in a vacuum. χ is the difference between the permeability limit and the vacuum boundary, and is called the electron convergence energy. F – output energy and it Fermi level F and the amount of energy equal to the vacuum surface difference. E_T – the limit of the external photoeffect, that is, the photon energy of the lowest energy that an external photoeffect can occur.

For the state where the Fermi level lies within the prohibited zone, i.e., the external photoeffect condition in a private semiconductor:

$$E_T = \chi + E_g \quad (13.37)$$

In a non-transition semiconductor, a phonon is also emitted:

$$E_T = \chi + E_g + E_p \quad (13.38)$$

Here E_p is the energy of the phonon.

Heavily Alloyed n - and p -Fermi levels in type semiconductors $\mathcal{E}_n$ and $\mathcal{E}_p$ The quasi-Fermi is divided into Fermi levels. For such cases, (13.38) the equation would be as follows:

$$E_T = \chi - \mathcal{E}_n \quad (13.39)$$

$$E_T = \chi + E_g + \mathcal{E}_p \quad (13.40)$$

The above-mentioned basic magnitudes for various semiconductor materials F – output energy, F - Fermi sathi, E_T – the values of parameters such as the external photo effect limit are listed in the table below.

Table 13.1

Output case for some semiconductor materials
and external photo effect border

Semiconductor	E_g , eV	F , eV	E_{th} , eV	E_{td} , eV	$\chi = E_g$, eV
AlSb (110)	1,5	4,86	5,15	5,75	3,65
Gas (110)	1,40	4,71	5,47	5,75	4,07
GaSb (110)	0,70	4,76	4,76	5,24	4,06
InAs (110)	0,41	4,90	5,31	5,58	4,90
InSb (110)	0,18	4,75	4,77	5,26	4,59
InP (110)	1,3	4,45	5,68	5,94	4,38
Yes (110)	1,09	4,83	5,10	5,45	4,01

Ge (111)	0,67	4,80	4,80	5,22	4,13
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Control Questions

1. What is internal photoeffect?
2. What is photoconductivity?
3. What is the relationship of photoconductivity to the recombination of charge carriers? and photoconductivity where diffusion is present.
4. What is the relationship of photoconductivity to charge carrier diffusion?
5. How is the Dember effect explained?
6. What is the photoelectromagnetic effect?
7. How to explain the photo-effect in the p-n transition?
8. How to explain the photo effect in the Schottky Barrier?
9. What is external photoeffect?

XIV bob.

OPTICAL AND PHOTOELECTRIC DEVICES

14.1. Photoelectric phenomena in p-n transition structures.

Photodiodes

A diode that depends on the feedback current is called a photodiode. P-n transition diodes with a reverse-directional voltage put on are usually used as photodiodes. When light quanta (photons) are absorbed in the p-n transition and in the adjacent semiconductor fields, new charge carriers (electron-hole pairs) are formed. Non-basic charge carriers that occur in regions that approach the p-n transition and whose width is less than the diffusion length L diffuse and are diffused towards the p-n transition, and under the influence of an electric field in it, they pass into the region in which they are main (Fig. 13.5). Because of this, the feedback current through the photodiode increases (Fig. 14.1). This current gain is called a photocurrent.

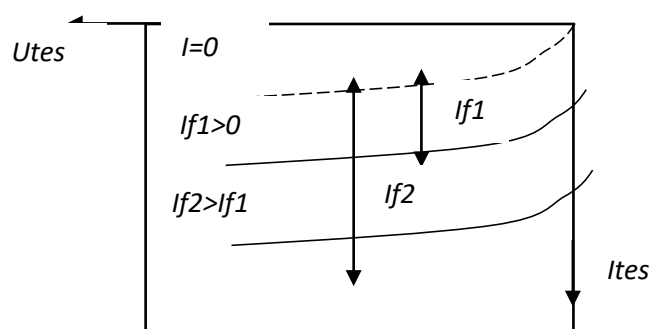


Fig. 14.1. Inverse Branch of Photodiode BAX in Different Light Flux

In the working range of inverse voltages, feedback currents (photocurrents) are practically not dependent on the voltage. The main characteristics of photodiodes include: The lux-ampere characteristic of a photodiode is that a photocurrent is bound to illumination, and the photocurrent is normally proportional to the illumination, i.e., $J_f \sim I$. From this it follows that the integral sensitivity of the photodiode does not depend on the given inverse voltage:

$$K = J_\phi / I$$

Photodiodes have a small inertia compared to photoresistors, that is, they are faster. In some photodiodes, the width of the base is 20–30 μm and the time of passage of the charge carrier through the base is $\tau_b \approx 50 \text{ ns}$.

But, p-n pass time

$$\tau_{p-n} \approx \delta / v_{\max},$$

This is the era δ – p-n-jump width, $v_{\max}$ is the greatest drift velocity of charge carriers, usually $\tau_{p-n} \approx 0.1$ value of ns.

The time constant of the photodiode is $\tau_{p-n} = r C_{p-n}$ depends on the voltage through the capacitance, and the photodiode base resistance depends on r , which has values in the order ns. The spectral characteristics of photodiodes are determined by large wavelengths – the width of the semiconductor's forbidden zone, the high absorption coefficient by small wavelengths, and the increased surface recombination. The maximum state of spectral characteristic is strongly dependent on the degree of growth of the absorption coefficient in a given semiconductor. For example, in a germanium photodiode, it determines the maximum position of the forbidden zone width ($\lambda_{\max} = 1.55 \text{ mkm}$) and does not depend on the base thickness. However, in silicon photodiodes, it is possible to shift the spectral characterization maximum in the range 0.6 to $1\mu\text{m}$ by varying the base thickness and surface recombination rate.

14.1.1. Calculation of photodiode photocurrent

To calculate the photodiode photocurrent, let's look at the photodiode described in Fig. 14.2. As shown in Fig. 14.2, the luminous flux is directed by the p-sphere along the p-n transition. Let the p-n transition be given an inverse voltage. Let's calculate the photocurrent for this case. Density of current passing through a photodiode

$$j = j_n + j_r + j_p = j_k + j_f \quad (14.1)$$

is defined by expression. where j_n is the current of electrons in the p-sphere: j_g is the current of electrons in the p-n transition region generated by thermal and photogeneration: j_p is the current of the holes in the n-sphere. On the other hand, the total density of the reverse current can be thought of as the sum of the unit current density j_k , and the photocurrent density j_f .

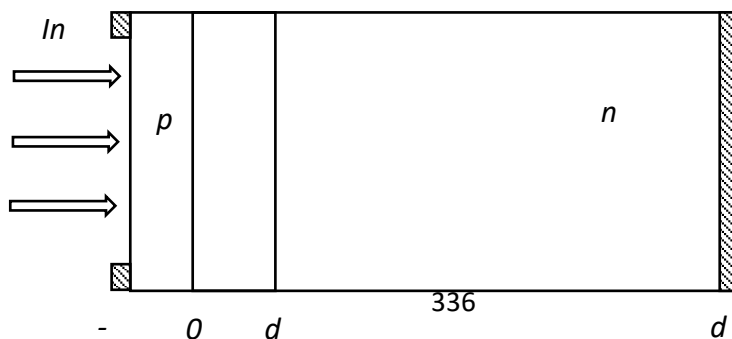


Fig. 14.2. Photodiode structure

The specific resistance of the p and n fields is small, in which the electric field strength is close to zero, not to mention the consequently drift current. If we call the flux of light entering an n-type semiconductor in the plane $x=0$ $I(0)$, then according to Bouger-Lambert's law, the flux of light inside x distance is

$$I(x) = I(0) \exp(-\alpha x) \quad (14.2)$$

So α is the coefficient of absorption of light. Absorption Light in Unit Size

$$dI(x)/dx = \alpha I(x)$$

is expressed by the number of photons absorbed in that volume,

$$\alpha I(x) = \hbar \omega \cdot q_1 \quad (14.3)$$

Putting the expression (13.3) in (1.81) is derived from:

$$\begin{aligned} G_n &= \eta_n q_1 = \eta_n \alpha I(x) / \hbar \omega = \eta_n [\alpha I(x) / \hbar \omega] \exp(-\alpha x) = G_{n0} \exp(-\alpha x) \\ G_p &= \eta_p q_1 = \eta_p \alpha I(x) / \hbar \omega = \eta_p [\alpha I(x) / \hbar \omega] \exp(-\alpha x) = G_{p0} \exp(-\alpha x) \\ G_{n0} &= \eta_n \alpha I(0) / \hbar \omega, \quad G_{p0} = \eta_p \alpha I(0) / \hbar \omega \end{aligned} \quad (14.4)$$

p-sphere $x=d_p$ relative to the luminous flux entering the plane

$$G_n = G_{n0} \exp[-\alpha(x+d_p)]$$

should be written.

1. The continuity equation for electrons in the p-field is as follows:

$$2. \quad \frac{d^2 n}{dx^2} - \frac{\Delta n}{L_n^2} + \frac{G_{n0}}{D_n} \exp[-\alpha(x+d_p)] = 0 \quad (14.5)$$

To solve this equation, we note the following boundary conditions:

$$\left. \frac{dn}{dx} \right|_{x=-d_p} = (s_p/D_n) \Delta n \Big|_{x=-d_p} \quad (14.6)$$

$$\Delta n(0) = n_p \quad (14.7)$$

s_p is the coefficient of superficial recombination.

$$j_n = \frac{e G_{n0} L_n}{a^2 L_n^2 - 1} \left\{ \frac{[D_n \operatorname{sh}(d_p/L_n) + s_p L_n \operatorname{ch}(d_p/L_n)] \exp(-\alpha d_p) - (a D_n + s_p) L_n}{D_n \operatorname{ch}(d_p/L_n) + s_p L_n \operatorname{sh}(d_p/L_n)} + \alpha d_n \exp(-\alpha d_p) \right\} - j_{sn}$$

In this case, $j_{sn} - p$ is the saturation current of a type thin-base diode. The rest of the jihads represent the electron carrier of the photocurrent, the j_{nf} .

Agar $d_p/L_n \ll 1$ va $\alpha d \ll 1$ bo'lsa, u holda

$$J_{nf} = - \frac{e G_{n0} d_p}{1 + s_p d_p / D_n} \left[1 - \frac{a^2 L_n^2 s_p d_p / D}{1 - a^2 L_n^2} \right]. \quad (14.8)$$

If the surface recombination rate is small ($s_p d_p / \ll 1$),

$$J_{nf} = -q_e C_{no} d_p, \quad (14.9)$$

that is, all the electrons that produce light in the p-field reach the p-n transition and contribute to photocurrent. It turns out that the smaller the s_p surface recombination coefficient, the higher the efficiency of the photodiode's work.

A large p-sphere thickness ($d_p/L_n \gg 1$) is not desirable, since less light reaches the n-sphere boundary, and the recombination in the p-sphere reduces the charge carriers generated by the light, the lower the photodiode efficiency.

2. The generative carrier of photocurrent in the p-n transition is soda;

$$J_{nf} = -eG_{p0} \int_0^d \exp[-a(x + d_p)] dx = eG_{p0} a^{-1} \cdot \quad (14.10)$$

$$\exp(-ad_p)[1 - \exp(-ad)] \quad (14.11)$$

From this expression it is clear that the smaller d_p , the greater the fold d (p-n transition thickness), the greater the J_{nf} current.

3. For j_{pf} hives via photodiode,

$$\frac{d^2 p}{dx^2} - \frac{\Delta p}{L_p^2} + (G_{p0}/D_p) \exp[-a(x + d_p)] = 0 \quad (14.12)$$

The continuum equation can be solved from the following boundary conditions

$$\Delta p(d) = p, \quad \Delta p|_{x \rightarrow 0} = 0 \quad (14.13)$$

Let's solve using the following:

$$j_p = [eG_{p0} L_p / (1 + aL)] \exp[-a(d + d_p)] - eD_p \overline{p_n} / L_p. \quad (14.14)$$

At the same time, he was the first person to do so.

$$j_{p\phi} = [eG_{p0} L_p / (1 + L_p)] \exp[-a(d + d_p)], \quad (14.15)$$

The second is the hive organizer of the dark reverse vine through the photodiode:

$$j_{ps} = [eD_p P_n / L_p]. \quad (14.16)$$

Usually, when photodiodes are made, the d_p width of the p-sphere is made very small, in which almost no light is absorbed. Chargers are not generated. That is why the organizer of the $j_{n\phi}$ is small and

$$j_\phi = j_{p\phi} + j_{ps} \approx eG_{p0} a^{-1} [1 - \exp(-ad)/(1 + aL_p)] \quad (14.17)$$

When the photodiode is given the correct voltage, the dark current j_{ns} and j_{ps} should be multiplied by $\exp(eV/kT)$, j_ϕ photocurrent in this case decreases due to the narrowing of the p-n transition area and the increase of $\exp(-ad)$. This means that when a large correct voltage is given, the full current of the photodiode becomes practically equal to the correct current in the dark, therefore, in this case, the diode will not be a photosensitive device.

14.1.2. Quantum efficiency and velocity of a photodiode

The ratio of the number of electron-pit pairs (j_{ϕ}/e generated by light to the number of photons emitted per unit ($I_0/\hbar\omega$) is called the quantum efficiency (efficiency) of the photodiode:

$$\eta_g = (j_{\phi}/e) / (I_0/\hbar\omega) = (I.R) \eta_f [1 - \exp(-ad) / (1 + aL_p)] \quad (14.18)$$

This is one of the main parameters that determines the quality of the photodiode. A parameter called photodiode sensitivity is often used, which is the ratio of photocurrent to light intensity (flux):

$$\beta_g = j_f/I_0 = e\eta_g / \hbar\omega = \lambda\eta_g / 1.24. \quad (14.19)$$

In this case, λ is expressed in μm , and β_g is expressed in A/W. The smallest repetition or maximum wavelength at which a significant photocurrent can be obtained is determined by the fact that the energy is equal to the width of the quantum semiconductor bandwidth.

$$\hbar\omega_{\min} = E_g \text{ from } \omega_{\min} = E_g / \hbar \text{ and } v_{\min} = E_g h,$$

$$\lambda_{\text{makc}} = hc/E_g \quad (14.20)$$

In the case of $\lambda > \lambda_{\text{makc}}$, a and G_0 are very small, there is no noticeable photocurrent, which means that light passes through the photodiode without forming electron-pit pairs.

In contrast, in the short wavelength ($\lambda < \lambda_{\text{makc}}$) region, the radiation is absorbed completely in the sensitive layer of the semiconductor. BUT,

$$ad_p \gg 1, \quad \exp(-ad_p) \ll 1$$

The main carriers of photocurrent, j_{gf} and j_{pf} , will be very small. The maximum (maximum value) in the intermediate area should be observed.

Indeed, the maximum value of η_g for a silicon photodiode is $\lambda \approx 0.8-0.9 \mu\text{m}$, and for a germanium photodiode $\lambda \approx 1-1.6 \mu\text{m}$.

The velocity of the photodiode is determined by $t_{\text{uch}} = d/v_{dt}$ (where v_{dt} is the maximum drift velocity of the charge carriers in a reverse-voltage p-n transition strong electric field) and the recharge time of the p-n transition capacity is determined by $t_c = C_{p-n} \cdot r_{yu}$ (r_{yu} is the charge resistance). V_{dm} (300K) $\sim 10^7$ cm/s for Ge, Si, GaAs. Hence, it is possible to increase the photodiode velocity by reducing the p-n transition width and surface again. The photodiode velocity again depends on the time of diffusion of charge carriers through the p-n transition volumetric charge region ($\tau_g = d_p^2/D_n$), in which case the p -field width can be increased by reducing d_p . The photo-response time of photodiodes lies in the range of $10^{-8}-10^{11}$ s (T=300K). Photodiodes can amplify periodically variable light streams, (optical signals) up to that interval, and then modify them.

There are several types of photodiodes: p-n and p-i-n- photodiodes,

- Shotki obstacle photodiodes,
- heterotransitive photodiodes,
- Avalanche photodiodes,

- Injecting photodiods.

14.2. Phototransistors. Photothyristors

A transistor that is excited by a luminous flux and can amplify the photocurrent at the same time is called a phototransistor.

Consider a double-pole phototransistor. The structure schemes of phototransistors are different. The phototransistor can be illuminated from different angles, it can be connected according to different schemes. We'll see one of those. If a voltage is given between the base and the collector, the collector transition is connected in the opposite direction, the emitter is not connected to the output circuit, then the circuit of the dual polar transistor becomes similar to that of the photodiode.

Unbalanced pairs of charge carriers are formed when photons are absorbed in the base and collector fields.

Non-basic charge carriers (holes at the n-base, electrons in the p-collector) are diffused towards the collector, the electric field pulling them into the collector passage and creating a photocurrent as it passes through. However, double-polar phototransistor is generally used in a common emitter circuit. So let's look at the double-pole phototransistor operation connected in this scheme.

Initially, we assume that the base electrode is not connected to the circuit ($I_b = 0$). In this case, the non-basic charge carriers pass through the collector p-n pass and generate a photocurrent. The electrons that produce light at the n-base and come from the collector are located in the potential pit, due to their accumulation, I_f both the emitter and collector current increase. Thus, the non-equilibrium principal charge carriers stored at the base of the double-polar phototransistor amplify the photocurrent, i.e., when illuminated, the collector current becomes $I_f = h_{21} I_p$, in which the coefficient is $h_{21} > 1$.

When the base electrode is connected to the outer circuit, the base current can be changed due to the influence of lighting. The degree of change depends on the base circuit resistance, which occurs as a result of the exit of non-neutral electrons into the outer base circuit, which reduces the electrons accumulated at the base, while the photocurrent gain is also reduced.

This means that the double polar phototransistor is most sensitive to the illumination of its field in a common emitter circuit, when the base is a ring. And the connection of the base makes it possible to control the phototransistor not only optically but also electrically, balancing the edge external effects. A thyristor, in which the connection voltage decreases with increasing illumination, is called a photothyristor.

One of the base areas of the photothyristor is illuminated (Fig. 13.10).

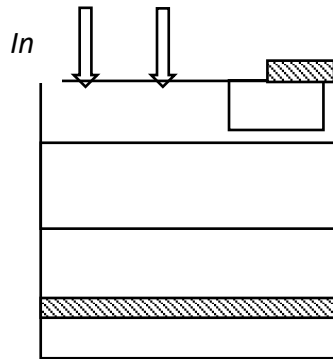


Fig. 14.3. Structure of the photothyristor

The reconnection of the photothyristor from the closed state to the open state occurs when the transfer coefficient of the thyristor structure over the current is multiplied by 1, as in a normal thyristor. Editor's Choice All Thus, the photothyristor is a controlled thyristor analogue, and its transition from the closed to the open state is controlled by a light impulse.

14.3. Photovaricans and Coordinate Sensitive Photoinstruments

14.3.1. Photovarikanlar

Special diodes (with p-n junction or MDYA structure) intended for use as capacities are called varicaps. The P-n transition varicap capacity can be expressed as:

U_k The difference in contact potentials when the external voltage is not put: A degree exponent for a constant sharp p-n transition is $k=1/2$. Photovoltaics are such devices that they convert light energy into an electrical signal and parameterize it. Its principle of operation is based on the dependence of the barrier capacity on light. Photo EYUK affects capacity as if it were an external voltage. The amount of capacity per unit surface without external voltage is $0.03 \mu\text{F}/\text{cm}^2$ for Si and $0.40 \mu\text{F}/\text{cm}^2$ for GaAs. The photosensitivity coefficient of photosensitivity represents the relative change in capacity due to the unit light flux falling on the surface. Its amount is $5,7 \text{ mW}^{-1}$, for Si structures and 240 mW^{-1} for GaAs. In the photovoltaic mode of operation, the photosensitivity is higher than in the photodiode mode. The temperature coefficient of the capacity is $1,5 \cdot 10^{-3} \text{ }^\circ\text{C}^{-1}$, for Si, and $0,8 \cdot 10^{-3} \text{ }^\circ\text{C}^{-1}$ for GaAs. Photovoltaic diodes, or MDYAs, built on the basis of structures, are considered promising because their capacity is tightly tied to voltage.

14.3.2. Coordinates – sezgir fotoasboblar

The output signal depends on the location of the illuminated part, called coordinate – sensitive photodevices. They are used as instruments to determine the coordinates in which the emitting objects are located in space, or as instruments to measure deformation.

A longitudinal photoeffect-based coordinate-sensitive photodetector is shown in Figure 14.4 and has a p-n transition consisting of both low-resistance p-layers and high-resistance n-layers. Phototension is obtained from the 1 and 2 electrodes in the n-layer. The principle of operation of such a photo acceptor can be explained by means of an equivalent chain shown in Fig. 14.4. The resistances of the layers p – and –n are P_n , P_p , and the thicknesses of w_n , w_p are related as follows: $P_n \gg P_p$, $w_n \ll w_p$. In this case, the p-layer can be considered as an equivalent surface relative to the n-layer. The p-n junction and the n-base can be thought of as diodes and resistors distributed over length. Photo with a valve in the illuminated part of the photodiode *E.Yu.K.* is formed, *E.Yu.K.* Under the influence, a current flows through the circuit of each of the elements, consisting of a diode and a resistor, and a voltage drops occur in each resistor. The U_{12} voltage in the 1st and 2nd electrodes is the sum of the voltages in the individual resistors. As can be seen in Fig. 14.4. $U_{12} = i_2 R_2 + i_3 (R_2 + R_3) + \dots + i_m (R_2 \dots R_m) - i_1 R_1$. Nur R_2 and R_3 When the beam moves between resistors R_2 va R_3 , $i_2 R_2$ becomes negative and the voltage U_{12} decreases. When the light reaches the middle of the sample ($x=0$), the sum of positive and negative values is equal and $U_{12}=0$. When the beam is moved further away (to the 2nd electrode side), the gesture of U_{12} changes to the opposite. (Fig. 14.4). The sensitivity of such devices is 1-1.5 mW/mW·mm.

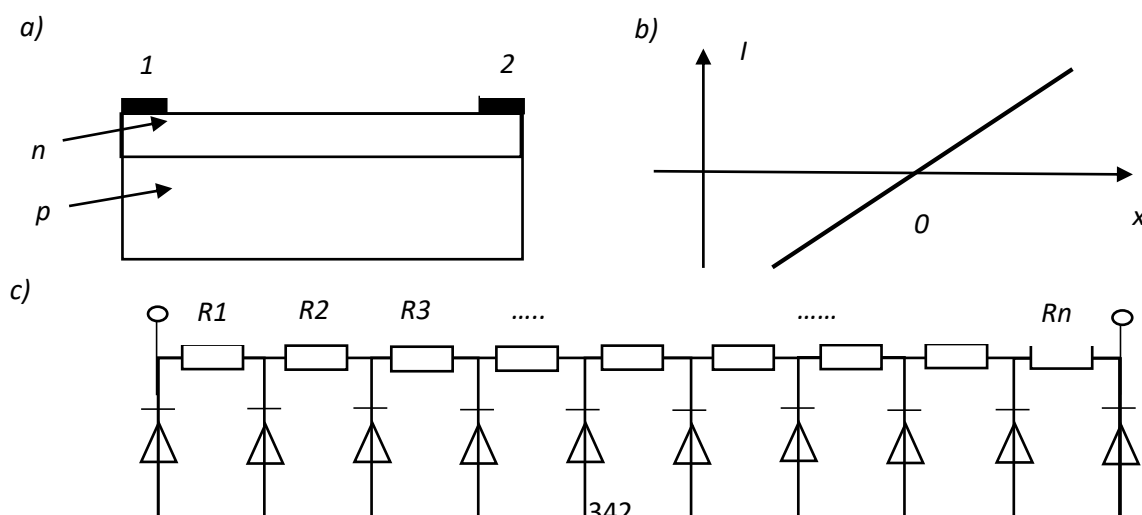


Fig. 14.4. The principle of operation of the photo acceptor

14.4. Solar cells and batteries

For five billion years, thermonuclear reactions inside the sun have been releasing large amounts of energy due to the transformation of a hydrogen nucleus into a helium nucleus. The temperature inside the sun is estimated to be between 15 and 20 million degrees. With 600 million tons of hydrogen being converted to helium every second, the mass of $2 \cdot 10^{27}$ tons of the sun has been reduced over billions of years by fractions of just one percent. The mass of the Sun is 330,000 times greater than that of Earth, and its diameter is $1.39 \cdot 10^9$ meters. From the sun, a beam of light with a power of about 1000 W per square meter of the Earth's surface falls constantly. In the spectrum of sunlight, energy is very unevenly distributed, with spectra outside the atmosphere and on the Earth's surface differing from each other.

Currently, various ways of using solar energy for human need have been developed. Along with the use of sunlight for heating buildings and water, drying fruits and other products, semiconductor photovoltaic elements have been created. Thanks to them, solar energy can be obtained directly electrical energy in semiconductor crystals due to an internal photoeffect. In this paragraph, we will look at semiconductor photocells and the process by which they convert light energy into electrical energy.

A semiconductor diode that converts solar energy into electrical energy is called a solar cell or photocell. It works without external stress. In a p-n transition photocell, the additional charge carriers formed by the action of light are separated by the p-n transition electric field, and the p-n transition electrons are collected in the n-field while the chain is far away. If an external circuit is connected, a current directed from the p-field to the n-field begins to pass through the loading resistor. Consequently, the full current through the p-n transition will R_{io} consist of a distinction of the direct current V in relation to the direct voltage and the photocurrent J_ϕ in the opposite direction:

$$J = J_s [\exp(aV/kT) - 1] - J_\phi, \quad (14.21)$$

From this:

$$V = \frac{kT}{e} \ln\left(\frac{J_\phi + J}{J_s}\right). \quad (14.22)$$

With the chain ring ($I=0$, $R_{io}=\infty$) V voltage p-n in transition *photo* $E.Yu.K.$ is equal to:

$$V = \frac{kT}{e} V_R = (kT/e) \ln \left(1 + \frac{J_\phi + J_R}{J_s} \right). \quad (14.23)$$

In the case of a short circuit ($R_{io}=0$, $V=0$ va $-I=I_{k,T}=I_\phi$), this mode of p-n transition is called the ventilated photoelement mode.

Ideally, all radiation that falls on the photocell is completely absorbed and recombination is considered to be non-existent. Without it

$$J_\phi = eS \int_{h\nu=E_g}^{\infty} \frac{dn_\phi}{d(h\nu)} d(h\nu) \quad (14.24)$$

In this case, $dn_\phi / d(h\nu)$ 1 cm² is the number of photons falling at 1 s per and the unit energy range. The spectral distribution of the flux of photons coming from the Sun is known, and (14.24) can be calculated as the inequal.

Fig. 14.5 illustrates the VAXs of ideal and true photoelements.

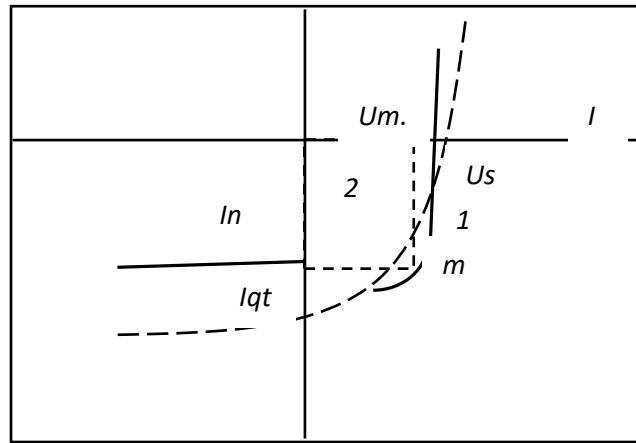


Fig. 14.5. Ideal (1) and Realistic (2) Photocells VAX i

The main parameter of a photoelement is its useful coefficient (F.I.K.), which expresses the ratio of the maximum power separated during the load to the falling radiation power:

$$F.I.K. = P_{max} / I_o \cdot S \quad (14.25)$$

(14.25) Express-based power dissociable in load

$$P = J \cdot V = J_s V [\exp(qV/kT) - 1] - J_\phi V \quad (14.26)$$

When $V=V_m$ the power has a maximum value: $P_{max}=I_m V_m$. Provided that the power has a maximum value, i.e. $(dP/dV)_{V=V_m} = 0$:

$$|J_M| = J_s (qV_m/kT) \exp(qV_m/kT). \quad (14.27)$$

With that condition, if we do some things,

$$V_m = V_{si} - (kT/q) \ln(1 + qV_m/kT). \quad (14.28)$$

Based on the above expressions (14.26) and (14.27)

$$P_{Makc} = J_\phi [V_{cio} - (kT/q) \ln(1 + qV_M/kT) - (kT/q)]. \quad (14.29)$$

As the bandwidth of the semiconductor's bandwidth E_g increases, J_s decreases, V_{si} increases, but J_f decreases. Hence, the P_{max} value depends on E_g (13.13).

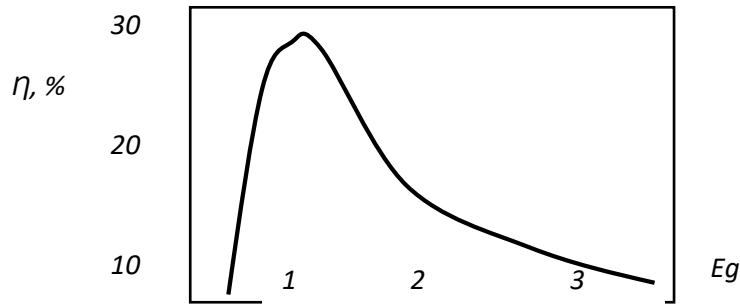


Fig. 14.6 Link of F.I.K. to E_g .

As can be seen from Fig. 14.6, it is possible to efficiently make a photocell from semiconductors with $E_g = 1-2$ eV.

14.4.1. Volt-ampere Characteristic of a Real Photocell

We have provided information about the ideal photoelement above. Taking into account the voltage drop (JR_s) in the p- and n-fields and contacts in real photoelements, (14.29) is written by modifying the expression as follows:

$$J = J_s \{ \exp[q(V - JR_s)/kT] - 1 \} - J_f \quad (14.30)$$

Expressing power:

$$P = \frac{kT}{e} \left(\frac{J + J_\phi}{J_s} + 1 \right) + JR_s \} J \quad (14.31)$$

Visibility is formed. Based on these statements, one can say the following considerations:

- 1) With respect to the ideal photocell, the current J increases more slowly when V increases in the true photocell.
- 2) When $V = 0$, $J_f \neq J_{qt}$;
- 3) In a true photoelement, the maximum power P_{max} (which is depicted in Fig. 14.6 as a striped rectangle, for an ideal FE it is equal to the face of the $V_m I_m$ rectangle) is much less;
- 4) If the recombination in the p-n transition is taken into account, the exponential representation of VAX to the expression of (14.28) is twice as small, which also reduces the rate of growth of the current, i.e., in which case I_s must be replaced by $I_i = en_i d / 2T$; Since it is $J_i \gg j_s$, this substitution also reduces V_{si} and P_{max} . The photocell is affected by electromagnetic radiation (solar energy is mainly collected in this range) in the wavelength range from 0.2 to 0.3 μm .

$$R_H = r_H np. \quad (14.32)$$

In which case the red limit of λ_{max} -absorption is, $N_f = I_0/h\nu$ is the density of the flux of photons in the unit spectral range, and $R_q(\lambda)$ is the repulsion coefficient.

As can be seen from the expression (14.32), in order to obtain a maximum photocurrent, it is necessary to reduce $R_q(\lambda)$ and multiply $\eta_i(\lambda)$ over the entire range.

The F.I.K of an ideal photoelement decreases as the temperature increases, the reason being that when the temperature increases, J_s or j_0 decreases, which decreases J_f , while in the dark, the current also decreases its growth rate due to exponential values of q/kT or $q/2kT$.

The more kata the illuminated face and the p-n transition face of the photoelement, the greater the force dissipated in the illumination. Unlike photodiodes, photovoltaic cells do not require speed.

14.4.2. Types of photoelements

Monocrystalline silicon-based photovoltaic cells are most commonly used and provide the conversion of solar energy directly into electric energy. Since photocells are expensive, their use on Earth is not yet very profitable from an economic point of view, and cheap polycrystalline silicon, amorphous silicon-based photocells would be widely used, silicon photo cells would make up no more than 12% of the F.I.K. Photovoltaic elements have also been developed. Calculations show that the height of the barrier is around 25% of the F.I.K. in the marginal where the height of the barrier is $\varphi = E_g$.

Preparation of photoelements based on polycrystals CdS, CdTe, Si GaAs is interesting. Their main virtue is their relative cheapness. However, their F.I.K. is low and their characteristics are slowly deteriorating.

Cheap amorphous silicon layers saturated with hydrogen are grown in glass bottoms 1-3 μm thick. *a-Si* Characteristics of E_g It is similar to that of a straight-band semiconductor with $= 1.6 \text{ eV}$. BUT, *a-Si* The F.I.K of the instruments made on the basis does not exceed 12%.

In photocells prepared on the basis of heterotransition, a wide-band semiconductor acts as an optical window for short-wave radiation. At these intervals, the spectral sensitivity of these photoelements is greater than that of homoetic photocells. Other measures can improve the efficiency of heterotransition photovoltaic cells. Solar cells are typically assembled from ordinary photocells to obtain large amounts of photoelectric energy. In solar cells, it is possible to obtain the necessary large voltage by connecting photovoltaic cells in series.

One way to get more photoenergy with less use of expensive semiconductors is to make "solar modules." Solar modules are a concentrator device that concentrates light over a large area into a small-sized (semiconductor) photocell. In

such a device, more photoelectric energy is obtained due to the fact that a large amount of light radiation falls on the photocell. Currently, solar modules have been developed for different shapes and sizes, as well as photocells made from different semiconductor materials.

14.5. Light-emitting diodes and lasers

14.5.1. Light-emitting diodes

A p-n transition diode that converts electrical energy into light energy is called a light-emitting diode. The conversion of such diode electrical energy into light energy is possible due to the radiant recombination of additional (surplus) charge carriers that occur under certain conditions in the p-n transition and in adjacent fields.

One convenient way to generate additional charge carriers is to conduct a straight current through a p-n crossing. In this case, non-equilibrium charge carriers are introduced (injected) into the p-n spheres, into the p-n transition region. The radiative recombination that occurs when a direct current passes through the P-n transition is called inductive electroluminescence. Possible views of radiant transitions are illustrated in Fig. 14.7.

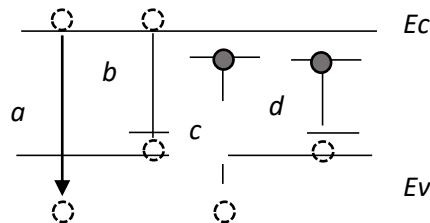


Fig. 14.7. Radiant recombination: a - inter-band electron and pit recombination; b - recombination of a free circuit with an accuser; C - recombination of free-hive with a donor; D - Recombination by Dependent Donor-Actor Pairs

Rate of interzone recombination (5.19, Fig. a)

$$R_H = r_H np. \quad (14.32)$$

In Balance:

$$R_H = r_H n_o p_o = y_n n_i^2. \quad (14.33)$$

From 14.32 and 14.33 comes the following expression:

$$R_H = R_o np / n_i. \quad (14.34)$$

Radiation recombination rates at the activator and donor centers, respectively, were (Fig. 14.7-b):

$$R_a = Y_a n N_a, R_d = Y_d p N_d. \quad (14.35)$$

At low T , the actuator-donor-dependent pairs contribute an important characteristic to radiation. In this case, the electron that was originally held by the donor is recombined with the electron that was originally held by the donor.

Under thermodynamic equilibrium, the rate of generation of electron-cavity pairs is equal to the rate of their radiation recombination. A relatively uncomplicated account is carried out as follows:

$$R_0 = \int_0^\infty r_0(v)dv = (8\pi/c^2) \int_0^\infty \frac{\bar{n}^2 a(v)v^2 dv}{\exp(hv/kT)-1}, \quad (14.36)$$

where n is the refractive index of light, and s is the velocity.

One of the main parameters of light diodes is F.I.K., which is measured in the ratio of fluorescent output power to a light diode:

$$\eta = N_f hv/JV, \quad (14.37)$$

N_f is the number of photons emitted in a unit time.

However, $J=qN_e$ and N_e are the number of electrons or holes that have passed through p-n in a unit time. Now

$$\eta = N_f hv/qN_e V \quad (14.38)$$

More might be added the concept of light emission coefficient (or optical efficiency):

$$\eta_{er} = N_f/N_{fi} \quad (14.39)$$

N_{fi} is the number of photons that are generated inside a light diode.

This magnitude

$$\eta_i = N_{fi} / N_e \quad (14.40)$$

The internal quantum represented efficiency. External quantum efficiency of a light diode:

$$\eta_t = N_f/N_e = \eta_i \eta_{er} = (F.I.K.) qV/hv \quad (14.41)$$

The values of η_t range from fractions of % to a few %.

The prohibited zone width of semiconductors used to make light diodes should be $E_g \geq 1.8 \text{ eV}$ (this is the red limit of visible light). One such semiconductor is gallium phosphide (GaP). It is a fault-band semiconductor in which there must be an intermediate input center in order for there to be effective radiation recombination. In GaP, this task can be performed by a nitrogen atom, which first captures the injected electron, then captures the swell, forming a bound exciton. Nitrogen-based GaP-based light diodes emit green radiation.

If Zn and O are simultaneously present in GaP, they form a neutral complex, in which case red ($h\nu=1.78 \text{ eV}$) radiation is observed. When Zn-O and N centers are introduced into GaP, both red and green radiation is obtained at the same time. GaAs_{1-x}P_x has found wide application in the preparation of light diodes that emit visible radiation in three-carrier solid solutions. Lead (x=0.65) and yellow (x=0.85)

radiation can be obtained if nitrogen is introduced into these fault zone compounds and GaP is grown on the bottom.

Light diodes manufactured on the basis of ZnS and SiC, emit ambient light, while the GaN-based p-n transition is located in the violet part of the radiation spectrum ($\lambda \approx 0.4 \mu m$).

The visible diodes are used in optoelectronic couples, both as indicators and displays.

For the preparation of diodes that emit infrared radiation, semiconductors with a $E_g < 1.5 eV$ are used (GaAs, $Ga_xIn_{1-x}As_{1-y}P_y$, $GaAs-Al_xGa_{1-x}As$). IQ diodes can be widely used as a source in fiber optic communication pathways.

14.5.2. Semiconductor lasers

The disadvantages of light diodes as a source are the relative size of the radiated power (in the range of 1mW) and the magnitude of the spectral range (100-500 A) of the radiation. These disadvantages are absent in injection lasers. In lasers, radiation occurs not as a result of spontaneous recombination of charge carriers, but as a result of stimulation or coercion. The word laser is a short term derived from the English name "Light Amplification by stimulated Emission radiation".

The conductivity electrons in a semiconductor are subject to the Fermi-Dirac distribution:

$$f(E, T) = [\exp(\frac{E - E_F}{kT}) + 1]^{-1}. \quad (14.42)$$

When the semiconductor is exposed to external energy, the distribution function of the electrons changes. However, the ratio of the probabilities that there are electrons in the two E_j and E_i states ($E_j > E_i$)

$$\frac{f(E_j)}{f(E_i)} = \frac{\exp(E_i - E_F)/kT + 1}{\exp(E_j - E_F)/kT + 1} \quad (14.43)$$

can be expressed in appearance.

As a result of the energetic action, if the distribution of electrons in the upper E_j state becomes greater than in the lower E_i state, i.e. $E_j > E_i$ and $f(E_j) > f(E_i)$ then

$$\frac{f(E_j)}{f(E_i)} = \frac{\exp(E_i - E_F)/kT + 1}{\exp(E_j - E_F)/kT + 1} \text{ or } \exp(\frac{E_j - E_i}{kT}) > 1.$$

For such a disparity to be fulfilled, there must be a $T < 0$. Therefore, in this case, the semiconductor is said to be in a negative absolute temperature state. The negative absolute temperature is not lower than the absolute zero temperature, but the effective statistical temperature corresponding to any pair of energy levels. If the probability of having electrons below the two energy levels is greater than the one above, in which case the temperature will be positive and vice versa. In

semiconductors, in order to establish a negative temperature state, an external energetic influence on the electron system must be exerted, transferring particles from the lower to the upper state.

This condition can occur in different ways:

- 1) injection of non-primary charge carriers from the P-n transition to fused semiconductors;
- 2) transfer of electrons to an excited state under the action of a strong electric field impulse to a single semiconductor;
- 3) in case of irregular circuit transition, their transfer to higher layers;
- 4) exposure to a semiconductor with accelerated electrons (energy in tens and hundreds, keV);
- 5) Semiconductor irradiation techniques with high intensity light are used.

The first way we see this laser is the injection laser –

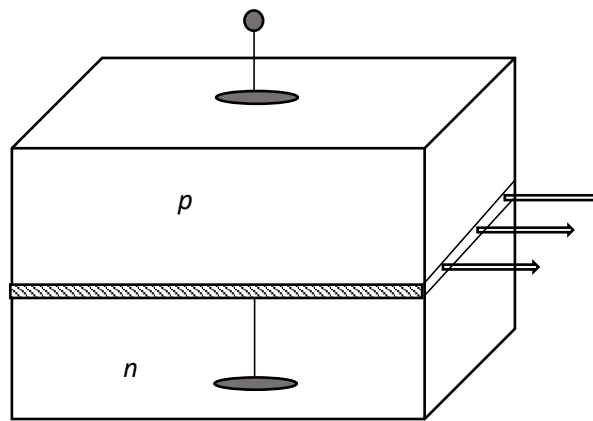


Fig. 14.8. P-N Slotted Injection Laser

We assume that both sides of the P-n transition are sober.

If the right voltage is given to such a p-n transition, the potential barrier to the p-n transition is reduced, while the densities of nonbasic charge carriers at its boundaries are increased. Consequently, electrons in the conduction zone begin to recombine intensively with spikes in the valent band, due to which the energy dissociated due to these transitions is emitted as a photon. For such an event to occur, it is necessary to have $V > E_g/e = \hbar\omega/e$. In this case, while the semiconductor laser is running, the potential barrier to the p-n transition is almost eliminated, and the tunneling effect does not occur.

The advantage of the injection laser is its continuous operation which can reach F.I.K. up to 100%. Its disadvantage is that the recombination area d is very narrow and its capacity is very small. The output capacity from recombination is very small. The repetition of the photon emitted by recombination is $\omega \approx E_g/\hbar$. If both sides of the p-n transition are in an equilibrium state, the energy $E_g + E_{Fp}(E_{Fn})$ is

needed to transfer the electron from the valent band to the permeability zone. The energy of the radiated photon is not enough for this. Photons that are emitted can therefore be emitted out. In this, in the presence of $E_{FC}-E_{F0}>\hbar\omega$, radiation stimulated by absorption prevails. The laser effect occurs only when the flux of photons (light beam) takes a certain direction. To do this, two opposite sides of the p-n transition crystal are polished. The light falling to this side intensifies and is emitted outward in the form of a thin handle.

This is because the phase, repeatability, and direction of the photons produced by the stimulated (forced) radiation generated by the insufficient external voltage current in the laser we are looking at will be homogeneous, i.e., the laser direction will be coherent. It is known that only coherent radiation can be interfered with and amplify each other.

It should also be noted that when a small correct voltage is given to the p-n transition, spontaneous and coherent radiation occurs in the laser. Only when the current reaches some J_{bo} 's value does the active field of the laser (p-n transition HZS and adjacent layers) emit a monochromatic and strongly directed light beam. Nowadays, injectable lasers are manufactured on the basis of smooth-band semiconductor compounds (GaAs, PbTe, PbSe, etc.) and solid solutions, which work in the coherent radiation range from 0.9 to μm . However, homotransition-based lasers can only work at the very large values of J_{bo} 's. For example, in the case of the passage of GaAs p-n, $J_{bo's}\equiv 5\cdot 10^4\text{A/sm}^2$. Strong heating of the diode does not allow the laser to work in a continuous marquee. They only perform so well at low temperatures.

Because of the above consideration, measures to reduce J_{bo} 's are prohibited. $\text{nAl}_x\text{Ga}_{1-x}\text{As-pGaAs}$ hetero structure injection lasers have some advantages over lasers only on the basis of homotransitions. For example, in them a value of is reached $J_{bo's}\leq 10^3\text{ A/sm}^2$, which makes it possible to work in continuous martyrdom at room temperature.

Control Questions

1. How do photoelectric phenomena shrink in p-n transition structures?
2. What is the working principle of photodiodes?
3. How is a photocurrent on a photodiode calculated?
4. How is the quantum efficacy of a photodiode explained?
5. How is the quantum efficacy of a photodiode explained?
6. What is the working principle of phototransistors?
7. What is the working principle of phototradars?
8. What is the working principle of photovaricans?
9. What is the working principle of coordinate sensitive photoinstruments?

10. What are the applications of coordinate-sensitive photodevices?
11. What are the structure and working principles of solar elements?
12. What are the structure and working principles of solar panels?
13. What is the volt-ampere characteristic of an ideal and true photoelement?
14. What are the types of photoelements?
15. What are light emitting diodes and what are their types?
16. What are the working principles of lasers?
17. What are the applications of lasers?
18. What are light emitting diodes and what are their types?
19. What is the working principle of semiconductor lasers?

APPENDICES

GLOSSARIY

ACCEPTOR is a defect in the lattice in which an unoccupied local line is present in the unoccupied state and in the waking state it attracts an electron from the valent field

AMORPHIC STATE is a solid non-crystalline state of matter; characterized by the isotrophy of properties and the absence of a melting point

AMORPH BODY is a Greek amorphous–formless–self-constituting body without the correct periodic arrangement

AMORPHIC SUBSTANCE is a solid substance without an ordered structure

ANISOLUTE BODY - a substance, the physical properties of which depend on the chosen direction in it

ANIZOTROOUS MEDIUM - a medium whose physical properties depend on the selected direction in it

ANIZOTROPATHY - non-uniformity of physical properties of the environment in different directions

MIXED SEMICONDUCTOR - A semiconductor with both ionic conductivity and conductivity properties inherent in electronic semiconductors at the same time

HOMOSEXUAL ENVIRONMENT - a medium in which certain physical properties do not depend on coordinates

HOMOGENOUS SYSTEM - a system with the same property in all parts

DIELECTRICS - substances that practically do not conduct electric current

DIODE is an electronic device with vacuum, semiconductor and gas charge, which conducts electric current in only one direction and has two circuits to connect it to an electrical circuit

DISLOCATION is a linear defect consisting in the violation of the arrangement of the atomic planes of a diclocatio–shift–crystal lattice

DONOR ACCESSION - An accessory in which atoms have donor properties

Defect of the donor-crystal lattice in which an electron is present at the local (local) level in the unawakened state, and in the awakened state, it gives an electron to the conductivity region

ELECTRIC TRANSITION - A sphere between quasi-homogeneous spheres in a semiconductor with different electrical conductivity or specific electrical conductivity (one of these quasi-homogeneous fields may also be metal)

ELECTRON CONDUCTIVITY - electrical conductivity caused by the movement of electrons

ELECTRON SEMICONDUCTOR - A semiconductor whose electrical conductivity (as opposed to an ion semiconductor) is determined by the migration of electrons

ELECTRON TRANSITION - A transition between two regions of an n-type semiconductor with different specific electrical conductivities

ELECTRON HIVE TRANSITION - the region in which the transition from cove conductivity to electron conductivity of a monocrystalline semiconductor occurs

ELECTRON CROSSOVER - A transition between two spheres of a semiconductor, one of which has n-type electrical conductivity and the other has p-type electrical conductivity

ELEMENTARY BRANCH - The part by which a crystalline atomic structure can be made into an entire crystal lattice by parallel motion in three dimensions

HETEROSTRUCTURA-A multiple-heterojunction semiconductor structure (structure)

HYBRID INTEGRATED MICROSTRUCTURE - an integral microstructure in which individual elements are formed independently constructively

HOMOTRANSITION is a contact of two spheres of different conductivity in a single semiconductor crystal as opposed to a heterotransition

IDEAL CRYSTAL - a crystal that does not have structural defects

ION SEMICONDUCTOR - A substance in which the energy range of its molecules between the electrolytic main band and the adjacent pole is greater than the electrolytic dissociation energy

ISOMORPHISM - Greek. Isos—the same, equal, and morphe-form, the complete similarity of the outer sides of its crystals with the atom-crystalline structure of substances with a similar chemical formula and identical chemical bonds

ISOTRO BODY is a body with the same properties in all directions

ISOTRO MEDIUM is a substance with the same physical properties in all selected directions

ISOTROPY - uniformity of physical properties of the environment in all directions

LOCATION DEFECT - a violation of the irregular exchange of states of atomic planes in domegranate volumetric centrifugal crystals or hexagonal superdense location crystals

CHEMICAL BONDING - A BOND FORMED BY THE AGGREGATION OF ELECTRONS ON DIFFERENT ATOMS OF A MOLECULE

INPUT SEMICONDUCTOR - A semiconductor having a donor and (or) acceptor input

A cluster is a system of a large number of atoms or molecules that are weakly bonded to each other

COMPENSATED FIELD – A semiconductor field with compensated semiconductor properties

COMPENSATED SEMICONDUCTOR - A semiconductor that has both donors and shareholders at the same time

HOLLOW SEMICONDUCTOR - Semiconductor consisting of hollows as the main charge carriers

HIVE FIELD - A field in which a pit in a semiconductor has electrical conductivity

KOVAK is an energetic state in which the electron is not charged in the valent field of a semiconductor

COVALENT BONDING - A chemical bond caused by the interaction of unpaired electrons having opposite directional spins occupying a single molecular orbit

COVALENT CRYSTALS - covalent chemical interatomic bonded crystals

CRYSTAL TEXTURE – Orientation of crystal grains in polycrystals

CRYSTALLINE STRUCTURE BASE - complete set of coordinates of centers of atoms in symmetrical independent spheres of a crystalline structure

CRYSTAL STRUCTURE - arrangement of atoms, ions, molecules in a crystal

CRYSTALLINE ASSOCIATIONS are monocrystalline spheres in which the crystal grid is located slightly differently from each other

CRYSTALLIZATION CENTER - the particle that causes the beginning of crystallization of a liquid

CRYSTALLIZATION TEMPERATURE is the temperature at which a phase transition from a liquid state to a crystalline state occurs

CRYSTALLIZATION - the formation of electrolytes in the solid state (amorphous state) electrolysis of vapors, solutions, alloys, as well as crystals during chemical reactions

CRYSTAL PHYSICS - a branch of solid state physics that studies the physical properties of crystals as well as the influence of external forces on the properties and structure of crystals

CRYSTAL SYMMETRY – the preservation of the self-alignment of crystals through twisting, rotation, parallel transfer or a combination of these operations

CRYSTAL SPECTROSCOPY - branch of spectroscopy, in which crystals study their various axis spectra in order to obtain information about their properties and structure

HEAD-SECTION IN CRYSTALS - 1) a plane in single-axial crystals that passes through a certain light beam and an axial axis; 2) a plane that passes through axial axes that intersect each other in two-axis crystals

DEFECTS IN CRYSTALS - defects in the crystal structure, a violation of the strict periodic arrangement of particles (atoms, molecules, ions) at the crystal lattice nodes

CRYSTALS - residual solids, arranged in a periodic order, forming particles

CRYSTALLITES are small crystals that have no distinct edges

CRYSTAL HEAD PLANE - **plane that** passes through the optical axis of a crystal and the direction of phase velocity of a light wave traveling through a crystal

AXIAL AXIS OF A CRYSTAL - the direction in which light propagates in a crystal without being refracted

CRYSTALLOGRAPHIC INDICES - three integers that determine the position of crystalline edges and atomic planes in space (Miller indices), as well as the direction of crystalline edges in relation to crystallographic axes (Weiss indices)

CRYSTALLOGRAPHY - crystals and *Greek. grapho - I write-* the science of the structure of crystals

CRYSTAL – *Greek. krystallos* – ice is a solid body with a three-dimensional periodic atomic or molecular structure and a regular symmetrical polygonal shape in equilibrium structure states

STRONGLY ALLOYED SEMICONDUCTOR - a semiconductor with very large concentration inlets (or structural defects of the crystal lattice) in which the force fields and wave functions of electrons located near them overlay each other due to the small distances between adjacent atoms

QUASI-STABLE PROCESS - a state in which the state within it does not change as a process spreads in a system

MICROSTRUCTURE - *Greek. micros-small and lot. Structuring - the structure* of a substance that is determined by means of a particle (axial or electron)

MICROSTRUCTURE - a microelectronic device having an equivalent assembly density and being considered a constructive whole, having an equivalent

assembly density in which at least five elements are located within the volume of 1 cm³ occupied by the structure

MIKRO is a Greek. *micro*– an additional phrase put before the name of a unit of physical magnitude to form the name of a unit of fraction equal to one millionth of a small initial unit

SUBSTANCE STRUCTURE - the collective name for microscopic and microscopic structure descriptions of a substance

MOLECULAR CRYSTAL - crystal located molecularly at the nodes of the crystal lattice

MONOCRYSTALS - a single crystal with a single crystal lattice

MONOCRYSTALLINE is a crystal having a single crystal lattice in its entire size

MONOMOLECULAR LAYER - a layer of substance equal to one molecule in thickness at the separation boundary of phases

COMPLEX SEMICONDUCTOR - A semiconductor, the main structure of which consists of atoms of two or more chemical elements

SUPERFINE STRUCTURE - the separation of atomic energy layers into nearby layers caused by the interaction of the magnetic field of atomic electrons with the magnetic moment of the atomic nucleus

OPTALATUM is a multilayer solid-state structure that affects electrons in addition to the periodic potential of the crystal lattice, as well as the artificial creation potential of the lattice constant

OPTICAL SMOOTH SURFACE - a surface in which the radius of curvature changes slightly relative to the wavelength of light across a surface

OPTOELECTRONICS is a field of electronics that deals with the conversion of electrical signals into optical signals and optical signals into electrical signals

ORGANIC SEMICONDUCTORS - solid organic substances having electronic or conductivity (or acquiring such properties under external influence) and a temperature positive coefficient of electrical conductivity

PARACRYSTALLINE - a molecular crystal in which crystalline and amorphous spheres are alternately located

PLANAR TRANSITION - A diffusion transition caused by the diffusion of an inlet through a crack in an absorber membrane laid on the surface of a semiconductor

P-N STRUCTURE - A STRUCTURE CONSISTING OF A P-SPHERE AND AN N-SPHERE BORDERING EACH OTHER

POLYCRYSTALS - solids composed of many crystallites (grains)

POLYMORPHIC TRANSITION TEMPERATURE - the equilibrium phase transition temperature of a substance at an invariable pressure from one crystal structure state to another

POLYMORPH TRANSITION - The phase transition of matter from one crystal modification to another

P-TYPE SEMICONDUCTOR - Semiconductor with spur conductivity

"EZOELCTRIC MATERIALS" - substances used in the manufacture of esoelectric converters, whose esoelectric properties are clearly expressed

THICK MEMBRANE INTEGRATED MICROSTRUCTURE - plated integrated microstructure whose membrane thickness is more than 1 μm

SOLID BODY is the aggregate state of matter; it is characterized by having a stable form and the thermal motion of the atoms that make it up in the form of small vibrations around their equilibrium states

RELAXATION TIME - The time it takes for a parameter of a system state to deviate from its equilibrium value by a factor

RELAXATION - lot. Relaxatio–Tension is a substitute process of thermodynamic equilibrium in a microscopic physical system

X-ray structural analysis - a method of studying the structure of a substance based on X-ray diffraction

FIXED TRANSITION - a non-contradictory transition that does not contradict the symmetry properties of the system

FIXED ZONE- An energy zone or collection of several overlapping energy zones formed by the displacement of one or more energy levels of individual atoms when atoms form a crystal

EFFECTIVE MASS - a magnitude about the size of a mass that describes the dynamic properties of a quasiaarra.

SURFACE WIDTH - The uncertainty of energy at the discrete energy level linked to the maximum lifetime of a quantum system at a given level

INVOLUNTARY LIVING on the surface is the time of residence at the level, which is determined by spontaneous radiation

RESIDENCE TIME AT A GIVEN LEVEL - *The time it takes for a system to reduce its residence time at a given level*

BANDWIDTH OF THE LAYER - the concentration of particles of a substance located at a certain energy level, that is, at a certain energy state

SURFACE PAIR - the states of a physical system corresponding to a certain energy level

SURFACE WIDTH - Uncertainty determined by the ratio of the magnitude of the Planck constant of energy of a given state to the residence time of that state

SMOOTH TRANSITION - A transition in which the width of the smooth change field of the inlet concentration is comparable to the width of the spatial charge field

SIMPLE SEMICONDUCTOR - A semiconductor whose main structure consists of atoms of a single chemical element

SPECIFIC DIELECTRIC LOSSES - a scalar magnitude that describes the distribution of dielectric losses in a dielectric volume and is equal to the ratio of dielectric losses in a volume of a dielectric to that volume

SORBATION - *Lat. sorbet – the absorption* of a liquid substance or gas (sorbate) from the environment of a swallowing liquid or solid body (sorbent)

LIQUID DIELECTRICS – molecular liquids with an electric conductivity $\sigma \leq 10^{-8} \text{ cm} \cdot \text{m}^{-1}$, electrons in molecules are covalently bonded, and between molecules there are Vander Vaals forces

LIQUID CRYSTALS are volatile states of certain substances with structural properties between solid crystals and liquids

FORBIDDEN ZONE is a realm of energy values that an electron in a crystal cannot possess

TEMPERATURE - a physical magnitude proportional to the average kinetic energy of the irregular motion of the particles that make up the system as well as characterizing the equilibrium state of the thermodynamic system

TENSORESIS EFFECT - A change in the electrical resistance of a semiconductor under the influence of mechanical shear

THERMOELECTRIC EFFICIENCY - the occurrence of electrical driving force when the contact temperatures are different in an electrical circuit consisting of different semiconductors or semiconductors and metals connected in series

THERMOELECTRIC DRIVING FORCE – an electrically motive force embodied in a thermoelectric effect

THERMOGALVANIC EFFECT - the occurrence of a transverse electric field strength in a semiconductor in the presence of a longitudinal temperature gradient and transverse magnetic field induction

THERMOGALVANIC MAGNETIC PHENOMENA - phenomena associated with the effect of a magnetic field on the electrical and thermal conductivity of solid conductors

THERMOPAIR - temperature recorder, the operation of which is based on the occurrence of thermal-electric driving force

THERMOMAGNETIC EFFECT - the appearance of a transverse temperature gradient in a semiconductor when the longitudinal temperature gradient and transverse magnetic field induction are induced

THERMORESISTOR (THERMORESISTOR) - A semiconductor resistance in which an active electrical resistance depends on temperature

INVERSE LATTICE - A point three-dimensional lattice specific to a particular crystal lattice, in which the distance between its nodes is inversely proportional to the distances between the corresponding nodes of the crystal lattice

FRICTIONLESS SEMICONDUCTORS - semiconductors with a bandwidth of exactly zero

FULL ZONE - permissible zone in which all energy levels at absolute zero temperature are banded with electrons

THOMSON ELECTROTHERMAL EFFECT - the dissipation or absorption of heat other than Joule heat when a current of a certain density passes through a homogeneous semiconductor and a longitudinal temperature gradient is present

TRANSLATION - *Lat. translatio-transmission, displacement-* the movement of an object at any distance along a line called the translation axis while keeping it parallel to itself in space

TUNNEL (-SI) EFFECT (-SI) - The phenomenon of particles passing through a potential barrier

AWAKENED ATOM - The state of an atom with a greater energy than in its ground state

AWAKENED STATE - A state of an energy greater than the smallest value of a discrete range of possible energies for an atom, molecule, or other quantum system

INDUCED CONDUCTIVITY - a change in the electrical conductivity of dielectrics and semiconductors under the influence of the current of particles (electrons, ions, etc.) whose energies are sufficient to form additional charge carriers

VACANCY - *lot. Vacans – Empty* – A crystal defect consisting in the absence of an atom or ion in a free crystal lattice node

VALENT ELECTRONS - The outer electrons of an atom

VALENT SPHERE - the domain of possible energy values of electrons in a crystal fully filled with valent electrons at absolute zero temperature

VALENT SPHERE - the highest of the filled spheres (the region of the greatest energies)

XEMOSORPTION - adsorption, which occurs in conjunction with the formation of chemical compounds

HALL'S COEFFICIENT (R) is $\vec{E} = R[\vec{j}\vec{B}]$ the coefficient of proportion in relation in $\vec{E}$ which is the transverse electric field strength, $\vec{j}$ - current density, $\vec{B}$ - magnetic induction

HALL EFFECT - $\vec{H}$ the formation of an electric field in a solid conductor in which a dense current $\vec{j}$ flows in a magnetic field $\vec{H}$ and $\vec{j}$ in a perpendicular direction

PRIVATE ELECTRICAL CONDUCTIVITY FIELD - a field of semiconductor with private conductivity properties

PARTIAL ELECTRICAL CONDUCTIVITY - the conductivity of an electric conductivity that occurs when a generation of pairs of electron conductivity of a semiconductor is generated in any way, such as heat, light, etc.

PRIVATE CONDUCTIVITY - Electrical conductivity of a chemical pure semiconductor

PRIVATE SEMICONDUCTOR - a semiconductor without donor and acceptor inputs

CLOSE ORDER is the relatively ordered arrangement of adjacent particles within small volumes of matter

SEMICONDUCTOR DEVICES - A common name for a variety of instruments whose performance is based on the properties of semiconductors whose performance has homosexual and non-heterosexual p-n transitions, heterojunctions

SEMICONDUCTOR DIODE - An electrical converter device that has an electrical circuit(s) and two outputs

SEMICONDUCTOR INTEGRATED CIRCUIT - an integrated circuit, the elements of which are created on the volume and (or) surface of the semiconductor material

SEMICONDUCTOR LASER - A laser created on the basis of a semiconductor crystal

SEMICONDUCTOR MATERIALS - A set of substances that clearly demonstrate semiconductor properties over a wide temperature range, including room ($T \sim 300$ K) and are widely used in the manufacture of semiconductor devices

SEMICONDUCTOR ALLOYING - Insertion of ingresses or structural defects into semiconductors in order to change their electrical properties

SEMICONDUCTOR - A substance that is between a conductor and a dielectric in its specific electrical conductivity and differs from a conductor in that its specific electrical conductivity is strongly bound to the temperature and concentration of the inputs

DENSE LOCATION – The location in which the maximum number of bubbles are located in a given volume

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Basic physical constants

Electron charge $qe = 4.8 \cdot 10^{-10}$ one. CGSE = $1,602 \cdot 10^{-19}$ Cl

Erkin electron mass $me = 9.11 \cdot 10^{-28}$ G

Velocity of light in vacuum $c = 2,998 \cdot 10^{10}$ cm/s

Boron radius $ao = 5.29 \cdot 10^{-9}$ cm

Plank doimiysi $h = 6.62 \cdot 10^{-34}$ J · s = $4.5 \cdot 10^{-15}$ eV · s

Dirak doimiysi $\hbar = h/2\pi = 1,054 \cdot 10^{-34}$ J · c $7,1 \cdot 10^{-16}$ эВ · c

Boltsman doimiysi $k = 1.380 \times 10^{-34}$ J/K = 8.62×10^{-5} eV/K

Heat energy $kT =$

= 25.9 MeV at room temperature

= 6.7 MeV at the temperature of dairy nitrogen

= 0.36 MeV at milk helium temperature

Amount of energy, 1 eV = $1,602 \cdot 10^{-19}$ Joules

The wavelength of light corresponding to the photon energy in vacuum λ_o (1 eV) = $1,239 \times 10^{-10}$ cm

The corresponding wave number for 1 eV, ν_o (1 eV) = $8.06 \cdot 10^{-4}$ cm⁻¹

Table 1.

Electrophysical Parameters of Pure Germanium and Silicon

Parameter	Digital Value	
	German	Silicon
1	2	3
Panjara doimiysi a , nm	0.565754 ± 0.00001 (298 K)	0.543072 ± 0.00001 (298 K)
Density (103 Kg/m ³)	5.3234 ± 0.0025 (298 K)	2.3283 ± 0.0003 (298 K)
Coefficient of Linear Expansion of Temperature	+0.07 (40 K)+1.5 (80 K)+2.20 (100 K)+5.75 (300 K)	0.05 (40 K)-0.77 (80 K)-0.31 (100 K)+2.33 (300 K)
Elastic coefficient: k ₁₁ k ₄₄ k ₁₂	12,886,7054,825	16,567,9536,386

Melting temperature, °C	937.2±0.5	1417±4
Melting heat	3290 ±90	5000±80 km
1	2	3
Evaporative heat	3719th ±900	4400±500
Specific heat capacity	25000 (300 K)11100 (80 K)	4.65 (300 C)1.24 (80 C)
Thermal conductivity	0.7 (300 K)3.0 (200 K)5.5 (50 K)11 (20 K) 10 (10 K) 0.5 (2 K)	1.3 (300 K)2.5 (200 K)13 (50 K)17 (20 K) 12 (10 K) 0.3 (2 K)
Debye harorati	353 (80 K)406 (300 K)	539 (80 K)689 (300K)
Dielectric absorption	16.3±0.2 (290 K)15.3±0.2 (77 K va 1 MHz)	11.7±0.2 (0-1 MHz)13.7 (24000 MHz)
Refractive index, <i>n</i>	4.01 (300K)	3.42 (299 C)
Outgoing work	7,52±0 ,16	7,68±0,16
Ban-zone width	1.0 56 (295 C)1.26 (0 C)	1.74 (295 C)1.93 (0 C)
Effective mass <i>m</i> * (in mo units): a) Heavy swells b) Light swells c) Electrons are the "longitudinal" constituent of the tensor <i>m</i> * d) The "transverse" constituent	0,340,0431,64 ±0,03 0,082 ±0,001	0,520,16 0,90 ±0,02 0,90 ±0,02

Table 2.

Basic physical parameters of semiconductor materials with elementary content

Mode (element)	Melting temperature, °C	Ban-zone width	Xarakatchanlik	
			$\mu_n, \text{m}^2/(\text{V}\cdot\text{s})$ (300° C)	$m_p, \text{m}^2/(\text{V}\cdot\text{s})$ (300° C)
In	2000	2,4		1
C (grafit)	4000	0,16	-	-
G (elms)	-	7,7	1800	1200
Yes	1420	1,78	1200	500
Ge	936	1,15	3900	1900
Sn (kulrang)	-	0,128	2000	1000
P (sariq)	44	3,35	-	-
R (qora)	-	0,53	200	350
Se (kristal)	220	0,92	-	5
They	452	0,56	1700	1200

Table 3.

Basic physical parameters of semiconductors with complex compounds

Mode (element)	Melting temperature, °C	Width of the restricted zone, eV	Mobility	
			$\mu_n, \text{m}^2/(\text{V}\cdot\text{s})$ (300° C)	$m_p, \text{m}^2/(\text{V}\cdot\text{s})$ (300° C)
Ag ₂ TeInSb				
H				
gTeMg ₂ SnH				
gSeInAs				
PbSe PbTe				
PbS InP	955 523 670		4000 80000	
GaAs	778 690 942		17000 300	
IgSb	1065 917 1114	0,27 0,29 0,32 0,32	100003000015	- 4000 100250 - 200
Mg ₂ Ge	1050 1280	0,48 0,46 0,88 0,96	00 2100	1500 800 800 700
GaSb	546 1115 725	0,96 2,1 2,2 0,96	6003500 4000	400 350 100 1000
	1102 1228 842	1,04 1,12 1,20	60 500 5000	70 80 - 100 150
MgSi	1045 1060	1,28 1.6 2.4 2,42	400 20 50	10100 -
Mg ₃ Sb ₂	2700 1232	2,4 2,5 5,1	450 50 60 -	
Ag ₂ S	1,975		200	
CdTe				
AlSb				
SiC				
Cu ₂ O				
ZnO				

Table 4.**Activation energies of access centers in silicon**

Periodic Groups	Element	Donor (D) or Actor(A)	D	A
I	Read	D	0,53	-
	At	D	0,86	-
		And	-	0,56
II	Zn	1A	-	0,50
		2A	-	0,88
III	In	And	-	0,0 72
	TO	And	-	0,091
	Ga	And	-	0,104
	In	And	-	0,246
V	R	D	0,0 63	-
	As	D	0,0 79	-
	Sb	D		

Table 5.**Activation Energy of Access Centers in Germany**

Embedded element symbol	Access item type, (A or D)	D	A
Read	D	0,016	-
	1A	-	0,064
With	2A	-	0,40
	D	0,84	-
At	1A	-	0,25
	2A	-	0,32
Zn	A	-	0,46
B	A	-	0,0166
TO	A	-	0,164
Ga	A	-	0,172
In	A	-	0,178
T	A	-	0,022

P	D	0,0193	-
As	D	0,0203	-
Sb	D	0,0155	-

Table 6.

Basic physical parameters of semiconductors with complex compounds

	Sb	D	0,0 155	-
	Mrs.	D	0,192	-
VII	Mn	A	-	0,244
		D	0,56	-
VIII	Fe	2A	-	0,63
		1A	-	0,56

	Co	D	0,495	-
		A	-	0,40
	NL	D	0,48	-
		A	-	0,35
	Pf	1A	-	0,0 64
		2A	-	0,736

Table 7.

Main properties of germanium and silicon

Physical Properties	Ge	Yes
Melting temperature, °C / Density, g/cm ³	936 / 5,323	1420 / 2,328
Coefficient of thermal expansion (25 °), 1 / grad	$6,1 \cdot 10^{-6}$	$4,2 \cdot 10^{-6}$
Heat conductive (25o) cal/s×cm×grad	0,14	0,20

Specific heat capacity, (0 – 100o), kal/s×sm×grad	0,074	0,181
Atom and'irligi	72,60	28,08
Panjara doimiysi, sm	$5,657 \cdot 10^{-8}$	$5,431 \cdot 10^{-8}$
Number of atoms (at 1 cm3)	$4,42 \cdot 10^{22}$	$5,22 \cdot 10^{22}$
Dimensional compression, cm2/din	$1,3 \cdot 10^{-12}$	$0,98 \cdot 10^{12}$
Dielectric constant	16	12
Activation Energy of Special Atoms, 96	0,75	1,1
Introductory Activation Energy of Atoms, 56	~0,01	~0,04
Special resistance (300° K) Om×sm.	47	3000
Electron mobility (300° K) cm2/V×s	3600	1260
Mobility of hives (300° K) cm ² /V×s	1700	500

Table 8.

Some parameters of semiconductors of type ^{AIII}B^V

Accumulation	Ter ,°C	ΔE_g, eV (300° C)	μ_n, m²/(V·s) (300° C)	m_p, m²/(V·s) (300° C)
Ge	937	0,67	0,39	0,19
Yes	1417	1,12	0,13	0,048
GaP	1450	2,25	0,011	0,0075

InP	1062	1,29	0,46	0,15
GaAs	1237	1,43	0,85	0,043
InAs	942	0,36	3,3	0,046
InSb	525	0,17	7,7	0,08

Table 9.

Some parameters of semiconductors of type $A^{III}B^V$

Material	Ban zone		Boundary wavelength, μm	Lattice constant a_0 , nm	Effective mass	
	Energiya, eV	Tour**			$\frac{m_e}{m_0}$	$\frac{m_h}{m_0}$
Ge	0,66	I	1,88	0,5657	0,22	0,30
Yes	1,11	I	1,15	0,5431	0,97	0,50
AlP	2,45	I	0,52	0,5451	-	0,70
AlAs	2,16	I	0,57	0,5661	0,15	0,79
AlSb	1,58	I	0,75	0,6135	0,12	0,98
GaP	2,26	I	0,55	0,5451	0,82	0,60
GaAs	1,42	D	0,87	0,5653	0,07	0,48
GASb	0,73	D	1,70	0,6096	0,04	0,44
InP	1,35	D	0,92	0,5870	0,08	0,64
InAS	0,36	D	3,50	0,6058	0,02	0,40
InSb	0,17	D	7,30	0,6479	0,014	0,40

Table 10.

Properties of Ge, Si and GaAs (300 K)

Property	Ge	Yes	GaAs
----------	----	-----	------

Amount of atoms p 1 cm ³	$4,42 \cdot 10^{22}$	$5,0 \cdot 10^{22}$	$2,21 \cdot 10^{22}$
Atomic mass	72,6	28,08	144,63
Drilling area, V/cm	$\sim 10^5$	$\sim 3 \cdot 10^5$	$\sim 4 \cdot 10^5$
Crystal structure	Elms	Elms	Spirituality
Zichlik, g/sm ³	5,3267	2,328	5,32
Relative dielectric absorption			
Effective state density in the permeability zone N_c , sm ⁻³	16	11,8	10,9
Effective state density in the valent zone N_c , sm ⁻³	$1,04 \cdot 10^{19}$	$2,8 \cdot 10^{19}$	$7,0 \cdot 10^{17}$
Effective mass m^*/m ; For Electron	$6,1 \cdot 10^{18}$	$1,02 \cdot 10^{19}$	$7,0 \cdot 10^{18}$
For the hive	M1.59	M 0.92	0,068
	mt0,082	mt0,19	
	ML*0.04	ppl*0.16	mRI*0.12
Electron approximation, eV	mpt*0.30	mpt*0,50	mpt*0.45
Forbidden zone width (300 K), eV	4,0	4,05	4,07
	0,66	1,11	1,43
Private concentration, sm ⁻³			
Panlara doimiysi, nm	$2,5 \cdot 10^{13}$	$1,6 \cdot 10^{10}$	$1,1 \cdot 10^7$
Linear expansion temperature coefficient, $\Delta L/L \cdot \Delta T$, °C ⁻¹	0,565748	0,543086	056534
Melting temperature, °C	$5,8 \cdot 10^{-6}$	$2,6 \cdot 10^{-6}$	$2,9 \cdot 10^{-6}$
Charge Carrier Residence Time, s	937	1420	1238
Drake's behavior, sm ² / (V·s): μ_n for electrons	10^{-3}	$2,5 \cdot 10^{-3}$	10^{-8}
For hives μ_p	3900	1500	8500
Phonon Roman energy, eV	1900	600	400
Specific heat, J/(g·°C)	0,037	0,063	0,035
Thermal conductivity (300 K), W/(cm·°C)	0,31	0,7	0,35
Thermal diffusion coefficient, cm ² /s	0,64	1,45	0,46
	0,36	0,9	0,44
Bug' bosimi, Pa	0.133 (1270 °C)	0.133 (1600 °C)	$1.01 \cdot 10^5$ (1050 °C)
	$1,33 \cdot 10^{-6}$ (800 °C)	$1,33 \cdot 10^{-6}$ (930 °C)	$1,01 \cdot 10^{-7}$ (1220 °C)
Outgoing case, eV	4,4	4,8	4,7

Jadval 11.

Output Case for External Photoeffect Boundary Energy

Semiconductor	E_g , eV	F , eV	Et , eV	Ent , eV	$\chi = Et - E_g$, eV
AlSb (110)	1,5	4,86	5,15	5,75	3,65
Gas (110)	1,40	4,71	5,47	5,75	4,07
GaSb (110)	0,70	4,76	4,76	5,24	4,06
InAs (110)	0,41	4,90	5,31	5,58	4,90
InSb (110)	0,18	4,75	4,77	5,26	4,59
InP (110)	1,3	4,45	5,68	5,94	4,38
Yes (110)	1,09	4,83	5,10	5,45	4,01
Ge (111)	0,67	4,80	4,80	5,22	4,13

Table 12.

Main properties of semiconductor materials

<i>In Mode</i>	<i>(Z) item sequence number, for (Excellent) connection</i>	<i>Lattice constant, ·10⁻¹ nm</i>	<i>Ter , ° C</i>	<i>Density, g/sm³</i>	<i>E_g, eV</i>	<i>μ_n , m²/(V · s)</i>	<i>μ_p , m²/(V · s)</i>
Elementary semiconductors							
Yes	14	5,43	1417	2,33	1,12	0,14	0,05
Ge	32	5,66	937	5,33	0,66	0,39	0,19
Herself	34	4,36	219	4,80	02,0	-	2 · 10 ⁻⁵
In	52	-	452	6,25	0,33	0,19	0,08
Compounds (AIII VV)							
BN	12	3,62	3000	3,49	6,0	-	-
To P	14	5,46	2000	2,37	2,45	0,008	0,003
G a P	23	5,45	1467	4,13	2,24	0,0019	0,0012
In P	32	5,87	1058	5,87	1,26	0,45	0,015
Al As	23	5,66	1770	3,60	2,16	0,028	-
G a As	32	5,65	1238	5,32	1,428	0,95	0,045
In Ash	41	6,06	942	5,67	0,356	3,3	0,046
Al S b	32	6,14	1069	4,28	1,58	0,02	0,055
G a S b	41	6,10	712	5,61	0,78	0,4	0,14
In S b	50	6,48	525	5,75	0,18	7,8	0,075

<i>In Mode</i>	<i>(Z) item sequence number, for (Excellent) connection</i>	<i>Panjara doimiysi, ·10⁻¹ nm</i>	<i>Melting temperatu re, T_{er}, ° C</i>	<i>Density, g/sm³</i>	<i>E_g, eV</i>	<i>μ_n , m²/(V · s)</i>	<i>μ_r , m²/(V · s)</i>
Compounds (All BVI)							
H g S	48	5,84	1480	7,73	1,78	0,07	-
Z n S e	32	5,66	1520	5,42	2,73	0,026	0,0015
H g Se	57	6,08	790	8,258	0,12	1,85	-
Zn T e	41	6,10	1239	6,34	2,23	0,053	0,003
Cd T e	50	6,48	1041	5,86	1,51	0,12	0,006
Hg T e	66	6,46	670	8,09	0,08	2,5	0,02
Other compounds							
If C	7	4,36	2800	3,21	3,12	0,033 – 0,1	-
P b S	49	5,94	1114	7,61	0,41	0,06	0,07
P b Se	58	6,12	1076	8,15	0,29	0,12	0,1
P b T e	67	6,46	917	8,16	0,32	0,18	0,09

Table 13.

Physical Properties of Certain Semiconductor Materials

<i>In Mode</i>	<i>Melting temperature, T_{er}, ° C</i>	<i>Lattice constant, Å</i>	<i>Density g /sm³</i>	<i>Boiling heat, With Kkal / mol</i>	<i>Thermal conductivity , kal/mol</i>	<i>Linear thermal expansion, $\alpha \cdot 10^4$</i>	<i>Width of the forbidden zone, Eg, eV</i>	<i>Electron Mobility μ_n, m²/(V · s)</i>	<i>Hives mobility, μ_p, m²/(V · s)</i>
Yes	1417	5,43090	2,32830	11,950	0,28	2,4	1,107	1600	500
Ge	937	5,65748	5,32600	8,100	0,14	5,75	0,66	3900	1900
In	452	a=4,495 b=3,74 c=5,912 d=2,86	6,25	4,180	0,03	Fe27,2 lc-1,6	0,33(0 K)	1890	790
Herself	219	a=b=4,36 c=4,96	4,80	1,500	0,02	47	1,79	—	1,0
Accumulations (AIII BV)									
GaAs	1238	5,65340	5,31700	25,18	0,37	6,0	1,428	8000	400
InAs	942	6,0584	5,667	18,4	0,29	5,19	0,356	22600	200
GaSb	712	6,09593	5,6137	15,56	0,105	6,9	0,78	4000	650
InSb	525	6,4794	5,775	11,4	0,15	5,04	0,18	78000	750
Gap	1467	5,4505	4,13	29,2	0,18	5,3	2,24	8500	435
InP	1058	5,8688	4,787	17,3	0,16	4,6	1,26	4500	150
Compounds (A ^{II} BVI)									
ZnS	1830	5,4109	4,089	-	0,006	6,2	3,6	140	5(700 K)
«	-	a=3,8200 c=6,2576	4,1	-	-	-	3,67	-	-
CdS	-	5,83	-	-	-	-	-	350	15

«	1475	a=4,136 c=6,749	4,820	-	0,048	-	2,42	295	-
β-CdSe	1239	6,084	-	10,5	0,0103	-	1,8	650	-
«	1258	a=4,299 c=7,020	5,81	-	-	-	1,74	600	-
HgSe	798	6,074	8,26	-	0,004	-	-0,07	(14-80)·10 ³	-
ZnSe	1820	5,6687 a=4,01 c=6,54	5,26 - -	- 15,6 -	0,046 - -	1,9 - -	2,70 - -	530 - -	28 - -
ZnTe	1290	6,089 a=4,27 c=6,99	5,72 - -	- - -	0,043 - -	8,3 - -	2,2 - -	300 - -	100 - -
Compounds (AVIV BVI)									
PhS	1113	5,9362	7,61	8,70	0,022	26	0,41	660	600
SnSe	875	a=3,99 b=4,32 c=11,18	5,080	7,55	-	-	1,08	-	90
PbS	1080	6,1243	8,15	8,5	0,020	10,4	0,29	1050	1000
SnSe	880	a=4,46 b=11,57 c=4,19	6,18	-	-	aa=-26 ab=35 ac=26	0,90		180
PhTe	924	6,45	8,16	7,5	0,028	19,8	0,32	1700	840
SnSe	806	6,3272	6,45	10,8	-	21	0,18	-	400
Compounds (AVIV BIV)									
a-SiC	~2800	a=3,0807 c=15,1174	3,2110	-	0,432	5,12	3,12	300	-
β-SiC	~2800	4,3596	3,210	-	0,432	3.8(200 °C)	2,2	100	-

Table 14.

Semiconductor Material Characteristics

Semiconductor	Width of the forbidden zone, (300K), eV	The energy of a pair of electron-hives	Atomic Number	Mobility (300K) $\mu, m^2/(V \cdot s)$		Residence Time, s	
				Electrons	Buckets	P-type electrons	N-type hives
Silicon	1.08	3.6	14	1500	500	$3 \cdot 10^{-3}$	$3 \cdot 10^{-3}$
German	0.67	2.94	32	3800	1800	10^{-3}	10^{-3}
Does not	6	10	6	1800	1200	-	-
Kadmiy selenidi	1.7	-	48.34	200	-	$10^{-2}-10^{-3}$	-
Cadmium Sulfide	2.4	7.3	4816	300	10	10^{-3}	$<10^{-8}$
Galliy arsenidae	1.43	6.3	31.33	8500	420	10^{-7}	10^{-7}
Galliy surmidi	0.67	-	31.51	4000	1400	10^{-8}	10^{-8}
Indium Phosphide	1.25	-	49.15	4800	150	-	-
Indiy Arsenidae	0.4	-	49.33	2000	240	-	-
Indy Surmidi	9.17	0.6	49.51	78000	750	10^{-7}	10^{-7}
Qo'rg'oshin surmidi	0.37	-	82.16	400	600	$2 \cdot 10^{-5}$	-

Table 15.

Semiconductor Material Characteristics

Group	Semiconductor	Band-zone width		Type of conductivity band absolute mimimics	$\left(\frac{dE_g}{dT}\right) \cdot 10^4$ (300 K), eV/K	$\left(\frac{dE_g}{dP}\right) \cdot 10$, eV/Pa	Effective mass		Refractive index, n	Relative electrical absorption	Grid constant a , $\text{\AA}$	Mobility	
		E_g (0K), eV	E_g (300K) eV				m_n^*/m_0	m_p^*/m_0				μ_n, cm^2 (V·c)	$\mu_{np} \text{ cm}^2$ (V·s)
1	2	3	4	5	6	7	8	9	10	11	12	13	14
IV	Yes	1,166	1,12	Regular 100	-2,3	-15	MM0.9	mpt0,56*	3,44	11,7	0,543	1350	480
	Ge	0,74	0,67	Regular 111	-3,7	5,0	2 mt0,19	mpl0.16* mpt0,33*	4,00	16,3	0,566	3900	1900
	α -Sn	-0,2		Regular 0000	-	5,0	M1,59 mt0,082 0,02	mpl0.04			0,6489	2000	1000
IV-IV SiC	A	3.0(6H)	2,8-3,2	Irregular	3,3				2,69 c		a0,30817 c1,51123 04359	400	
	β	2,68	2,2	Irregular					2,65⊥c	10,2			
	Herself	1,95	1,74	Regular 0001	-14	-20		0,12	5,56 c 3,72⊥c	8,5		1	
YOU	In	0,334	0.32	Regular 0001	-0,3	-19	0,038⊥	0,26	3,07 c 2,68⊥c	5,0 c 2,2⊥c	0,4538	1100	
	BP		2	Irregular				0,10⊥	2,6	6,9			

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