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Двухжидкостная модель сверхпроводимости в теории БКШ

Исследовано изменение поведения энергетического спектра коллективизированных валентных электронов при сверхпроводящем фазовом переходе. Коэффициенты канонической трансформации Боголюбова-Валатина выбраны в соответствии с двухжидкостной моделью сверхпроводимости. Новый "полупроводниковый" энергетический спектр получен как решение линеаризованного уравнения движений взаимодействующих электронов. Минимизацией термодинамического потенциала получена новая функция заполнения частиц f_k, которая отличается от принятой в теории БКШ. Энергетическое распределение нормальных и сверхпроводящих частиц получено в соответствии с двухжидкостной моделью сверхпроводимости. Представлена физическая картина корреляции между квазичастицами, формирующими Куперовские пары. Туннелирование нормальных квазичастиц в системах M-I-S и S_1-I-S_2 качественно исследовано с помощью двухжидкостной модели сверхпроводимости.

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Two-Fluid Model of the Superconductivity in the BCS's Theory

The coefficients of Bogolubov-Valatin's transformation are chosen in accordance with the two-fluid model of superconductivity. The energy spectrum of the superconducting quasi-particles is obtained as a solution of the linearized equation of motion of the interacting particles. The energy distribution of the superconducting and normal quasi-particle is discussed from a new viewpoint. The correlation between the quasi-particles forming the Cooper's pair is discussed in accordance with the proposed ideas. The tunnelling of the normal quasi-particles in systems M-I-S and S_1-I-S_2 is investigated qualitativly.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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1. Introduction

The superconducting phenomenon is a remarcable example of quantum effects on a macroscopic scale. In the superconductor a finite part of electrons is condensed in a "macromolecule", extended in the whole volume of the system, and capable to move as a whole. The superconducting state of metals can be described either by the microscopical theory of Bardeen-Cooper-Schriffer or by the macroscopical two-fluid theory of Ginsburg-Landau. Between them exists a certain mathematical correspondence. But in our opinion the physical correspondence between these theories is not satisfactory. The main purpose of the present work is to correct same misinterpretations and to present a full and perspicuous physical picture of the processes occurring in the superconductor. In order to understand and substantiate the new physical idea it is necessary to discuss many well-known mathematical outcomes and physical assumptions.

2. Mathematical Description of the Electron-Phonon System

In many-valence metals the high density of collective valence electrons (CVE) ensures a considerable average kinetic energy of the conduction quasi-electrons (CQEs). Therefore the crystal potential influence on the CQE motion is very small and the individual-particle model gives a fairly good description of normal metals. The essential change of CVE energy spectrum takes place only near the boundary of the Brillouin zones, along which the electron-lattice interaction is very strong and therefore the energy suffers abrubtion. In many-valence superconducting metals a large number of breaking off surfaces pass close to the Fermi surface. The real Fermi surface is similar to that, which is obtained in free-electron approximation [1]. Considering all this, one may assume that with a sufficient extent of accuracy the CVE energy spectrum in the normal state of a metal is described by the formula

$$\epsilon_k = \frac{\hbar^2 k^2}{2m^*}. \qquad (1)$$

The isotopic effect in superconductors [2,3] verifies Fronlich's assumption that the electron-phonon interaction favours energetically the superconducting state. In many-valence superconductors the electron-phonon interaction is strong because a large number of boundaries of Brilluoin zones pass close to the Fermi surface. Therefore for the subsequent description of the phase transition from normal to superconducting

state it is essential to investigate the behaviour of the system of CVE and phonons, considering only linear terms in their interaction. In this approximation and in the second-quantization representation the Fronlich's Hamiltonian can be written in the following form

$$H = \Sigma \epsilon_k a^+_{k\sigma} a_{k\sigma} + \frac{1}{2} V(q) a^+_{k+k\sigma} a^+_{k'+q\sigma} a_{k'\sigma'} a_{k\sigma} + \\ + \Sigma \hbar\omega_q (b^+_q b_q + \frac{1}{2}) + \Sigma g_q a^+_{k+q\sigma} a_{k\sigma} (b_q + b^+_{-q}). \qquad (2)$$

Here $a^+_{k\sigma}$, $a_{k\sigma}$, b^+_q and b_q denote the creation and the destruction operators of the CVE and the phonons, correspondingly. The last term in (2) describes the electron-phonon interaction, which in the first approximation determines the electrical resistivity of metals, and in a second approximation, makes it possible to find the CVE self-energy and to determine the dynamic correlation between the CQEs, which comes from the virtual exchange of phonons (one electron polarizes the lattice, and the other one interacts with it [5, p. 181]). The parameter of the electron-phonon interaction g_q is determined by the equality

$$g_q = i \frac{4\pi e^2 q}{q^2 + \gamma^2} \sqrt{\frac{N\hbar Z}{2M\Omega_p}}. \qquad (3)$$

In (3) N - denotes the CVE density, γ^{-1} - the screening length ($\gamma^2 = \frac{6\pi Ne^2}{\epsilon}$) and $\Omega_p = \sqrt{\frac{4\pi NZe^2}{M}}$ - the ion plasma frequency.

In metals the elastic waves are considered in the harmonic approximation. The phonon scattering and the simultaneous creation $b^+_{q_1} b^+_{q_2}$ and destruction $b_{q_1} b_{q_2}$ of two phonons are neglected. Therefore in the calculation of the matrix elements of b_q and b^+_{-q} only the emission or the absorption of one phonon may be taken into account. Thus in order to obtain the effective interaction between CQEs, it is convenient to take the matrix elements of the phonon operators while the Fermi operators remain in working order /5, p.181/ In this approach, taking into account the existence of zero oscillations only, the matrix elements of the operators b_q and b^+_{-q} are determined by the following formula /5, p.181/

$$\langle 0|b_q|1\rangle = -g_q \sum \frac{a^+_{\kappa'-q\sigma'} a_{\kappa'\sigma'}}{\epsilon_{\kappa'} - \epsilon_{\kappa'-q} - \hbar\omega_q}, \qquad (4a)$$

$$\langle 1|b^+_{-q}|0\rangle = -g^*_{-q} \sum \frac{a^+_{\kappa\sigma'} a_{\kappa+q\sigma'}}{\epsilon_{\kappa} - \epsilon_{\kappa+q} - \hbar\omega_q}. \qquad (4b)$$

Taking into account that $g^*_{-q} = g_q$, and substituting $\kappa = \kappa' - q$ in (4b) and (4) in (2) one obtains

$$H = \sum \epsilon_k a^+_{k\sigma} a_{k\sigma} + \sum \hbar\omega_q \left(b^+_q b_q + \frac{1}{2}\right) +$$

$$+ \frac{1}{2} \sum \left| V(q) + \frac{4\hbar\omega_q g_q g^*_q}{(\epsilon_k - \epsilon_{k-q})^2 - (\hbar\omega_q)^2} \right| a^+_{k+q\sigma} a^+_{k'-q\sigma'} a_{k'\sigma'} a_{k\sigma}. \qquad (5)$$

From (5) it is seen that for CQEs with energies ϵ_k and ϵ_{k-q}, which satisfy the inequality $|\epsilon_k - \epsilon_{k-q}| < \hbar\omega_q$, the dynamic attractive interaction, which comes from virtual exchange of phonons with frequency ω_q, decreases the screenin repulsive Coulomb interaction. Taking into account both interactions, an energy interval is to be supposed to exist, in which the dynamic attraction between CQEs is stronger than the Coulomb repulsion. The magnitude of this energy interval depends on the density and the energy spectrum of the CVE, and on the strength of the electron-phonon interaction.

3. Chioce of the Approximation

In the ground coherent state the quasi-momentum is conserved and the maximum number of negative matrix elements is obtained under the condition that all pairs have the same wave vector $q = K + K'$. The conservation law of the quasi-momentum strongly restricts the number of virtually exchanged phonons in dependence on the q. In the absence of an external field the largest macroscopic number of allowed exchange processes, which lead to the maximum reduction of the total energy of CQEs, is obtained when they form Cooper pairs with $q = 0$ /7, p. 155/. Although the pairs with $q = 0$ have negligibly small statistic weight (of order $1/N$)m still just they bring the essential contribution to the interaction energy, because the coherent effects compensate the insufficient contribution from a statistical point of view /8, p.49/; /9, p.144/. Really, if the terms

with $q = 0$ are regarded on the basis of the perturbation theory, then, in some approximation, their contribution to the energy of the normal state of a large system is negligibly small, but for superconductivity just these terms are the most essential. Cooper has shown that the degenerate Fermi gas with attractive forces between is unstable under the formation of pairs of CQEs with opposite spins, and equal, but opposite quasimomentum /10/. Taking into account all these consideration, one may drop in /5/ sum taken over K' and σ', retaining only terms with $K'=-K$ and $\sigma'=-\sigma$. Thus the reduced Hamiltonian of Bardeen-Cooper-Schriffer is obtained /5, p.188/

$$H_{BCS} = \Sigma \epsilon_k (a_k^+ a_k + a_{-k}^+ a_{-k}) + \\ + \Sigma <-\kappa, \kappa | W | k, -k> a_\kappa^+ a_{-\kappa}^+ a_{-k} a_k . \qquad (6)$$

Since the dynamic attraction between CQEs comes from virtual exchange of phonons with frequency $\omega_q \leq \omega_D$ (where ω_D is the Debye frequency of the crystal lattice), then the Fermi distribution of CVE strongly restricts the number of CQEs participating in that exchange to the energy interval $|\epsilon_k - \mu| < \hbar\omega_D$. Therefore one may assume to a good approximation the following model potential

$$W_{\kappa k} = \begin{cases} -V & \text{in the region } |\epsilon_k - \mu| < \hbar\omega_D ; |\epsilon_\kappa - \mu| < \hbar\omega_D \\ 0 & \text{outside this region} \end{cases} \qquad (7)$$

In a superconductor the number of superconducting CQEs is not conserved, but the

average value of the total number of CVE is conserved. Therefore in order to get rid of the necessity to take into account explicit-ing the constancy of the total number of particles

$$N = \Sigma (a_k^+ a_k + a_{-k}^+ a_{-k}) \qquad (8)$$

in accordance with general laws of statistics /11,p.92/ it is necessary to minimize the Gibbs' free energy $\Omega = H - \mu N - TS$, which at $T = 0^\circ K$ has the form of the quasi-particle Hamiltonian /6,p.276/, /11,p.299/, /12,p.259/, /13,p.122/, /14,p.31/, /15, p.53/.

$$\tilde{H} = \Sigma \xi_k (a_k^+ a_k + a_{-k}^+ a_{-k}) + \Sigma W_{\kappa k} a_\kappa^+ a_{-\kappa}^+ a_{-k} a_k . \qquad (9)$$

In the (9) the notation of quasi-particle energy $\xi_k = \epsilon_k - \mu$ is accepted, where μ is the CVE chemical potential, determined from the condition for conservation of the average total number of CVE.

4. Diagonalization of the Hamiltonian

In the mean field approximation the quasi-particle Hamiltonian (9) can be diagonalized by means of Bogolubov /16/ -Valatin /17/ transformation

$$c_k^+ = u_k a_k^+ - v_k a_{-k}, \qquad (10)$$

$$c_{-k} = u_k a_{-k} + v_k a_k^+ .$$

The eigenvalues of the eigenvectors in (10) are determined by the system of equations of motion for a_k^+ and a_{-k} /5,p.196/,/9,p.141/

$$ih\dot{a}_k^+ = -\xi_k a_k^+ - \Sigma W_{\kappa k} a_\kappa^+ a_{-\kappa}^+ a_{-k},$$
$$ih\dot{a}_{-k} = \xi_k a_{-k} - \Sigma W_{k\kappa} a_k^+ a_{-\kappa} a_\kappa \qquad (11)$$

/15,p.64/ , after their linearization. The parameters u_k and v_k are real numbers (for ensuring the reality of the CQEs occupation numbers /11,p.298/), which are connected by the equation $u_k^2 + v_k^2 = 1$. Moreover u_k is an even and v_k is an odd function of the wave vector K /5,p.197/,/12,p.252/. The linearization of the system of equations of motion (11) is a simple generalization of the Hartree-Fock method and consists of replacement of the operator products $\hat{B}_\kappa = a_\kappa a_{-\kappa}$ and $\hat{B}_\kappa^+ = a_{-\kappa}^+ a_\kappa^+$ with their average values $B_\kappa^* = \langle N+1 | a_\kappa^+ a_{-\kappa}^+ | N-1 \rangle$ and $B_\kappa = \langle N-1 | a_{-\kappa} a_\kappa | N+1 \rangle$, correspondingly. This corresponds to decoupling of the chain of equations for the Green function /18/ in the first order approximation. Introducing the notation

$$\Delta_k = -\Sigma W_{k\kappa} B_\kappa \qquad (12)$$

the system of equations (11) can be written in the following linearized form

$$ih\dot{a}_k^+ = -\xi_k a_k^+ + \Delta_k^* a_{-k},$$
$$ih\dot{a}_{-k} = \xi_k a_{-k} + \Delta_k a_k^+ . \qquad (13)$$

The linearization (13) of the system of equation (11) physical corresponds to the replacement of the interaction between CQEs, coming from the virtual exchange of phonons, with the interaction between CQEs and the mean field (12) of the remaining CQEs. The system of equations (13) has solutions of exponential form $a_k^+(t) = a_k^+(0)\exp(i\lambda_k t/\hbar)$ and $a_{-k}(t) = a_{-k}(0)\exp(-i\lambda_k t/\hbar)$ under the condition that λ_k is a solution of the equation

$$(\lambda_k - \xi_k)(\lambda_k + \xi_k) = \Delta_k \Delta_k^* . \qquad (14)$$

In this way it is obtained that $\lambda_k = \pm\sqrt{\xi_k^2 + \Delta_k^2}$ /9, p. 141/. The sign of λ_k corresponds to the sign of ξ_k.

The results, described above are well known and are generally accepted. However in our opinion their physical interpretation is incorrect. Indeed, from the abive conclusion it follows that if ξ_k is a quasi-particle energy of CQEs and describes their energy spectrum in the normal state of the metal, then λ_k is the quasi-particle energy of CQEs, forming Cooper's hard correlated pairs in the superconducting state of the metal, and therefore it ought to describe their energy spectrum. In such a way, using the relation $E_k - \mu = \lambda_k$, the semiconductor model of the energy spectrum of superconducting CQEs is obtained which is described by the following formula /19/,/20/,/21, p.330/, /22, p. 571/,/23, p.447/

$$E_k = \mu \pm \sqrt{\xi_k^2 + \Delta_k^2} . \qquad (15)$$

In fig. 1 the energy spectra of the normal CQEs ξ_k and of the superconducting CQEs E_k

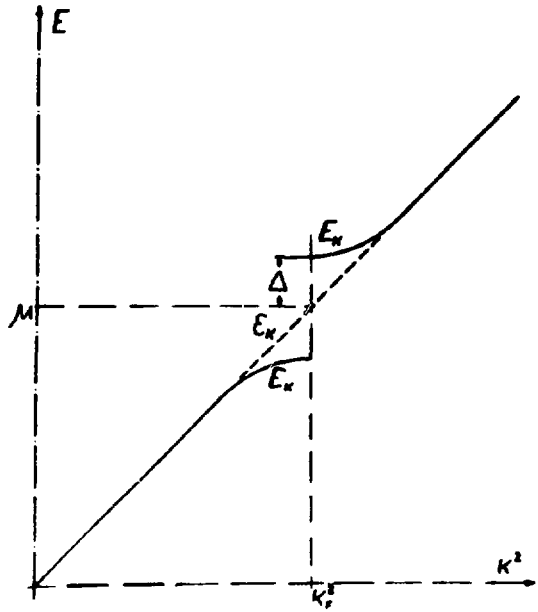

Fig. 1. The energy spectra of cooperons E_k and bogolons ϵ_k .

are presented. From the figure it is seen, that the width of the Cooper's energy gap between the spectra ϵ_k and E_k has its largest value on the Fermi level, where the number of the CQEs participating in the virtual exchange of phonons, has the maximum value.

5. The Energy Distribution of the CQEs

Since the operators c_k^+ and c_k act on the vacuum wave function $|\Psi_0\rangle$ ($c_k^+|\Psi_0\rangle = |\Psi_k\rangle$ and $c_k|\Psi_0\rangle = 0$), then for the determination of the CQE number ($\hat{n}_k = a_k^+ a_k$) and the condensation amplitude ($\hat{B}_k^+$ and $\hat{B}_k$) /1, p.117/ it is necessary to express the creation a_k^+ and destruction a_k operators of the interacting CQEs as functions of creation c_k^+ and destruction c_k operators of the weak interacting quasi-particles. In the BCS approximation the interacting between the new quasi-particles is neglected. So by means of a transformation, which is inverse to the transformation (10)

$$\begin{aligned} a_k^+ &= u_k c_k^+ + v_k c_{-k}, \\ a_{-k} &= u_k c_{-k} - v_k c_k^+ \end{aligned} \qquad (16)$$

after the regrouping of the operators one obtains

$$\hat{n}_k = u_k^2 c_k^+ c_k + v_k^2 c_{-k} c_{-k}^+ + u_k v_k (c_k^+ c_{-k}^+ + c_{-k} c_k), \qquad (17)$$

$$\hat{B}^+ = u_k^2 c_k^+ c_{-k}^+ + v_k^2 c_{-k} c_k + u_k v_k (c_{-k} c_{-k}^+ - c_k^+ c_k). \qquad (18)$$

Taking into account the definitions of the occupation number functions of the strong interacting CQEs n_k and of the weak interacting superconducting quasi-particles f_k,

and also of the condensation amplitude B_k from (17) and (18) we obtain the following equations:

$$n_k = u_k^2 f_k + v_k^2 (1 - f_{-k}), \tag{19}$$

$$B_k = u_k v_k (1 - f_{-k} - f_k), \tag{20}$$

If (20) and (12) are combined, one obtains an equation, which determines in a self-consistent way the equilibrium magnitude of the energy gap

$$\Lambda_k = -\Sigma W_{k\kappa} u_\kappa v_\kappa (1 - f_{-\kappa} - f_\kappa). \tag{21}$$

6. Comparison of the Ground State Energy of the Superconducting Models

Here it is necessary to correct some wrong physical interpretations of the mathematical results of BCS theory. Really, in the accepted interpretation of BCS theory all Cooper pairs have an equal energy, coinciding with the chemical potential of CVE, and therefore, all of them lie on the Fermi level [7, p. 159]. The excitations have the energy spectrum (15). Outside the Fermi sphere the quasiparticle excitations have the energy $E_k = \mu + \sqrt{\xi_k^2 + \Delta_k^2}$, and inside the Fermi sphere they have the energy $E_k = \mu - \sqrt{\xi_k^2 + \Delta_k^2}$. According to the BCS theory at $T = 0°K$ all

CQEs, which in the normal state of the metal occupy the states in the energy interval $\mu - \hbar\omega_D \leq \epsilon_k \leq \mu$, in the superconducting state of the metal are condensed in Cooper pairs. But the corrected calculation shows, that for raising the energy of CQEs from states with quasi-particle energy $\xi_k < 0$ to the Fermi level, an energy $E_S - E_N = 0.5N(\mu)(\hbar\omega_D)^2$ is necessary. This energy is $(\frac{\hbar\omega_D}{\Delta})^2$ times larger than the maximum energy gain of the superconducting state in comparison with the normal one. Since in fact the superconducting state has a lower energy, such an explanation of the results of the BCS theory is wrong. The condensation of the Cooper pairs, i.e., of the correlated quasi-particles, has to be understood in the sence that all of them have one and the same wave vector $q = K + K'$. But this is just the necessary condition, which secures the maximum number of negative matrix elements in the coherent ground-state. In the absence of external fields the coherent state with $q = 0$ has the lowest energy and as follows from (14) the quasi-particle energy of the correlated CQEs must be $E_k = \mu + \lambda_k$, while the quasi-particle energy of the uncorrelated CQEs is ϵ_k. From fig. 1 and formula (15) it is seen, that inside the Fermi sphere the energy branch E_k of the correlated CQEs lies below the energy branch ϵ_k of the uncorrelated CQEs, while outside the Fermi sphere the energy branch ϵ_k lies below the energy branch E_k. Therefore the quasi-particle distribution, which is described by the formula (19) is energetically more favourable. At $T=0°K$ inside the Fermi sphere the CQEs occupy the correlated states, but out-

side the Fermi sphere they occupy the uncorrelated states.

Besides the diagonalization (13) the above assumption may be motivated by showing that in the two-fluid model with the semiconducting energy spectrum of the correlated states, but with one-particle occupation in BCS's sence, is not energetically favourable. Really, let us assume that at $T = 0°K$ all the CQEs occupy only the states of the correlated quasi-particles with the energy spectrum E_k (15) and the energy distribution (19). Then the transition of the metal from a normal to a superconducting state changes the total energy of the CQEs by the amount

$$E(\Lambda) - E(0) = \frac{N(\mu)}{\sqrt{\mu}} \Big[\int_0^{\hbar\omega_D} \Big(1 - \frac{\xi}{\sqrt{\xi^2+\Lambda^2}}\Big) \sqrt{\xi^2+\Lambda^2}\ \sqrt{\mu+\xi}\, d\xi - \int_0^{\hbar\omega_D} \Big(1 + \frac{\xi}{\sqrt{\xi^2+\Lambda^2}}\Big) \sqrt{\xi^2+\Lambda^2}\ \sqrt{\mu-\xi}\, d\xi + 2 \int_0^{\hbar\omega_D} \sqrt{\mu-\xi}\, \xi\, d\xi \Big] . \quad (22)$$

In (22) the notation $N(\mu)$ is used for the density of one-electron states of the CQEs with a fixed spin orientation per unit energy at the Fermi surface. By means of a simple rearrangement from (22) one obtains

$$E(\Lambda) - E(0) = \frac{N(\mu)}{\sqrt{\mu}} \int_0^{\hbar\omega_D} (\sqrt{\mu+\xi} - \sqrt{\mu-\xi})(\sqrt{\xi^2+\Delta^2} - \xi)\, d\xi . \quad (23)$$

The equation (23) is exact. It shows, that in this model the phase transition of the metal from a normal to a superconducting state has to increase the total energy of the CQEs. This follows from the fact, that the density of states of the CVE increases with the increase of the energy. Therefore the energy density of states outside the Fermi sphere is larger than the energy density of states inside the Fermi sphere, while the change of the energy spectrum is symmetrical.

7. Physical Discussion of the Bogolubov-Valatin Transformation

Starting from (10) many physicists assume, that in a superconductor the quasi-particle is a linear combination of an electron and a hole. However this assumption is wrong [15, p.61]. Indeed, from the equations of motion (11) it follows, that in a superconductor the time evolution of the creation operator a_k^+ of the CQEs is connected first with the direct creation of a CQE, which is proportional to the quasi-particle energy ξ_k, and next with the creation of the CQE by means of the virtual exchange of phonons, which is described by the second term. Therefore in the second term there is one destruction operator a_{-k} and two creation operators $a_\kappa^+ a_{-\kappa}^+$. The definition of the condensation amplitude Δ_κ^* shows that it is an average of two creation operators, which create successively CQEs pair. Since in the accepted approximation (12) the virtual exchange of phonons is replaced by a mean field, then in the left-hand side of the equations there is only

one (creation or destruction) operator, while in the right-hand side of the equations stand both operators C_k^+ and C_{-k}. It will be more correct if an operator of the dynamical field $\hat{A}_k$ is introduced in consideration to describe the virtual exchange of phonons /15,p.66/. Then the Bogolubov-Valatin transformation (16) acquires the sence of operator equations, so that both terms on the right-hand side describe the appearance or annihilation of a quasi-particle, only if their second terms take into account the process of virtual exchange of phonons.

Thus the first term in the right-hand side of (16) takes into account the probability of the direct creation of CQE and therefore the first term in the right-hand side of (19) is connected with the occupation number function of the correlated superconducting CQEs f_k. The second term in (16) takes into account the probability of creation or destruction of CQE by means of virtual exchange of phonons and therefore the second term in (19) is proportional to the probability of the absence of superconducting CQEs $(1-f_{-k})$. From what has been said above it follows that in a superconductor the CQEs are distributed in both states /24, p. 32/. In the correlated states of cooperons $U_k^2 f_k$ with an energy spectrum E_k, and in the uncorrelated states of bogolons $U_k^2(1-f_{-k})$ with the energy spectrum ϵ_k of the normal CQEs. As far as we know, the excited CQEs in superconductors are called bogolons in BCS theory. A part of the correlated CQEs can be transferred in the uncorrelated states only expending an energy, which is necessary to overcome

the energy gap between both energy branches /15, p. 38/, /25/. In this way the Gorter-Casimir two-fluid model of superconductivity is naturally introduced /26/ . Many observed properties of the superconductors can be understood by means of the two-fluid model. The achievements of the Ginsburg-Landau's phenomenological theory /27, 28/ for description of the superconducting properties of metals prove that the application of the two-fluid model is valid.

8. Choice of Coefficients in the Bogolubov-Valatin Transformation

By means of the transformation (16), using the formulae (17) and (18), the BCS Hamiltonian (9) can be rewritten in the following form /11, p.299/

$$\tilde{H} = 2\Sigma \xi_k v_k^2 + \Sigma \xi_k (u_k^2 - v_k^2)(c_k^+ c_k + c_{-k}^+ c_{-k}) + \\ + 2\Sigma \xi_k u_k v_k (c_k^+ c_{-k}^+ + c_{-k} c_k) + \Sigma W_{\kappa k} \hat{B}_\kappa^+ \hat{B}_k , \qquad (24)$$

where the operator $\hat{B}_\kappa^+$ is determined from equation (18). In the Hamiltonian (24) only the terms, containing the products $f_k = c_k^+ c_k$ and $1 - \hat{f}_k = c_k c_k^+$ have diagonal matrix elements. Therefore, suing (24), the total quasi-particle energy of the CVE can be written in the form /11, p.299/

$$E = 2\Sigma \xi_k v_k^2 + \Sigma \xi_k (u_k^2 - v_k^2)(f_k + f_{-k}) +$$

$$+\Sigma W_{\kappa k} u_{\kappa} v_{\kappa} (1 - f_{\kappa} - f_{-\kappa}) u_k v_k (1 - f_k - f_k) . \quad (25)$$

By variation of (25) with respect to the parameter u_k and taking into account the relation $u_k^2 + v_k^2 = 1$, we get the following condition for minimum of the total quasiparticle energy (or the Gibbs' free energy at $T = 0^\circ K$) /11, p. 299/, /29, p. 193/

$$2\xi_k u_k v_k + (u_k^2 - v_k^2)\Sigma W_{kk} u_{\kappa} v_{\kappa} (1 - f_{\kappa} - f_{-\kappa}) = 0 . \quad (26)$$

Using the relation (21) one can rewrite the condition (26) in the following form /11, p. 300/

$$2\xi_k u_k v_k = (u_k^2 - v_k^2)\Delta_k . \quad (27)$$

Due to equation (26), which is the second relation between the parameters u_k and v_k, all the terms in the Hamiltonian (24) containing one pair of operators $c_k^+ c_{-k}^+$ disappear /29, p. 192/. In the first-order approximation of the perturbation theory they bring to divergent integral /11, p. 300/. Both the relations between the parameters u_k and v_k, the condition (27) and the relation $u_k^2 + v_k^2 = 1$ have a common solution, which can be written in the following analytic form:

$$u_k^2 = \frac{1}{2}\left(1 + \sqrt{1 - \frac{\Delta_k^2}{\xi_k^2 + \Delta_k^2}}\right) = \frac{1}{2}\left(1 + \frac{\xi_k}{\lambda_k}\right), \quad (28a)$$

$$v_k^2 = \frac{1}{2}\left(1 - \sqrt{1 - \frac{\Delta_k^2}{\xi_k^2 + \Delta_k^2}}\right) = \frac{1}{2}\left(1 - \frac{\xi_k}{\lambda_k}\right). \qquad (28b)$$

The expressions (28) are represented by double equalities in order to show how to select the signs, standing in front of λ_k and ξ_k [9, p. 142]. Moreover, it is natural, that if the second term in (16) takes into account the probability of the virtual exchange of phonons then the parameter v_k ought to be proportional to the value of the mean field parameter Δ_k. Really, after a simple rearrangement equation (28b) can be rewritten in the following form

$$2v_k^2 = \frac{\Delta_k^2}{\lambda_k^2 + \lambda_k \xi_k}. \qquad (28bb)$$

The equation (28bb) shows that the sign and the value of Δ_k determine the sign and the value of v_k. Thus we obtain that the choice of the sign in (28b) is physically correct. Using the relations (28), one can rewrite the condition (26) in the following form

$$2u_k v_k = \Delta_k / \lambda_k . \qquad (29)$$

As the parameter u_k has a constant sign, then the relation (29) shows that the sign of v_k is changed by changing both the signs (of Δ_k and of λ_k). Therefore the sign of v_k is changed by crossing the Fermi level and by changing the wave vector K. On fig. 2 the solid curve describes

the square of the coefficient u_k as a function of the energy ϵ_k of the normal CQEs. The dashed curve describes the square of the coefficient v_k in the above function.

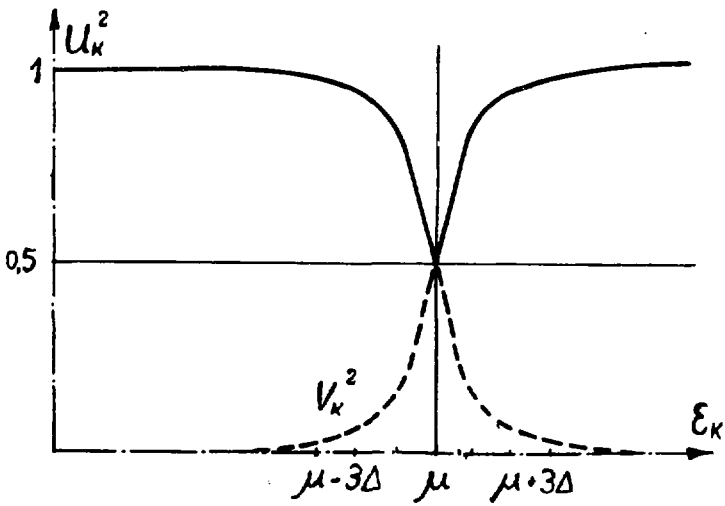

Fig. 2. Plots of the coefficients u_k^2 and v_k^2 as functions of the normal electron energy ϵ_k.

9. Self-Consistent Gap Equation

With the use of the above parameters of the transformation, the occupation number function of the cooperons f_k, minimizing the Gibbs' free energy $\Omega = H - \mu N - TS$, can be represented by the following formula

$$f_k = [\exp(\beta\lambda_k) + 1]^{-1}. \qquad (30)$$

In (30) we have accepted the notation $\beta^{-1} = K_0 T$, where K_0 is the Boltzman constant. Substituting in (19) the occupation number function f_k, given by expression (30), we get the following expression for the occupation number function n_k of CQEs /12, p.270/, /13, p.125/

$$n_k = \frac{1}{2}\left[1 - \frac{\xi_k}{\lambda_k}\,\mathrm{th}\left(\frac{\beta\lambda_k}{2}\right)\right]. \qquad (31)$$

Taking into account the conservation law of the total number of CVE we obtain that the chemical potential μ can be determined from the equation

$$\sqrt{\mu_0^3} = \sqrt{\mu^3} + \frac{3}{8}\,\frac{(\hbar\omega_D)^2}{\sqrt{\mu}}\left[1 - \sqrt{1+a^2} + a^2 \ln\left[\frac{1+\sqrt{1+a^2}}{a}\right]\right]. \qquad (32)$$

In (32) we denote $a = \frac{\Lambda}{\hbar\omega_D}$. Relation (32) shows that the change of the chemical potential is small and therefore may be neglected. Substituting (28) and (29) in (21) we obtain the following equation determining in a self-consistent way the value of the Cooper energy gap

$$\Delta_\kappa = -\sum W_{\kappa k} \frac{\Delta_k}{2\lambda_k} \operatorname{th}\left(\frac{\beta\lambda_k}{2}\right). \qquad (33)$$

The transition temperature T_c is defined as the boundary between the normal and the superconducting state of a metal. Therefore substituting $\Delta = 0$ in (33) we obtain an equation determining the transition temperature T_c. Since the energy gap between both energy spectra is increased when the temperature is decreased, the gap acquires its maximum value $\Delta_0 = 1.76 T_c$ at $T = 0°K$.

10. The Energy Distribution of Bogolons and Cooperons

Expressions (19), (20) and (28), together with fig. 3 show, that at $T = 0°K$ the CQEs distribution in a superconductor is the following: a part of CQEs occupy the superconducting states u_k^2 of cooperons with energy spectrum $E_k = \mu - \sqrt{\xi_k^2 + \Delta_k^2}$, which are found in a surface layer with thickness $\hbar\omega_D$ below the Fermi level. The remaining part of CQEs occupy the normal states v_k^2 of bogolons with energy spectrum ϵ_k, which are found above the Fermi level. Naturally, the CQEs, which are in states below the cooperons layer inside the Fermi sphere can not take part in the virtual exchange of phonon and therefore they have the energy spectrum of the normal CVE (1). This distribution follows from the fundamental assumption of the two-fluid model, applied to equation (19). Really, in the two-fluid model

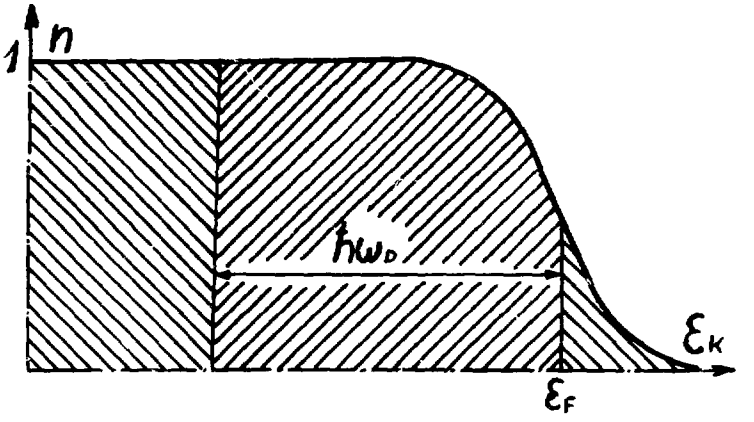

Fig. 3. The distribution of cooperons and bogolons in K-space at $T = 0°K$.

the CQEs are found in both states, in the normal and in the superconducting state[24, p. 32]. Hence it follows that if the state is not occupied by a cooperon, it must be occupied by a bogolon. Therefore the first term in (19) describes a cooperon distribution and the second term describes a bogolon distribution.

However, although the bogolons have the energy spectrum ϵ_k of the normal CQEs, their number of states is limited in a superconductor. As at $T = 0°K$ all bogolon states v_k^2 above the Fermi level are occupied by CQEs, bogolons must not take part in transport processes. At $T = 0°K$ all bogolons take part in the virtual exchange of phonons, giving rise to the energy spectrum E_k of the cooperons. Taking into account the energy spectra of CQEs in s superconductor, the

energy distribution of cooperons and bogolons are presented on the fig. 4.

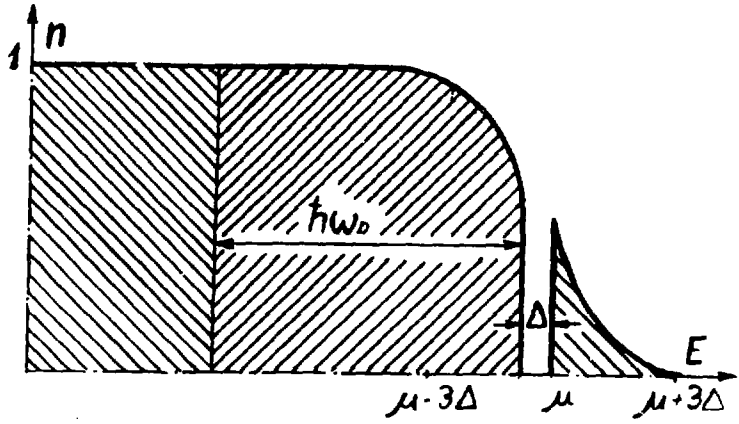

Fig. 4. Energy distribution of cooperons and bogolons at $T = 0°K$.

11. The Change of the Total Energy of CQEs and the Lattice at Superconducting Phase Transition

In order to prove that the assumed physical model of superconductivity is correct, it is necessary to compare the total quasiparticle energy (25) with the total quasiparticle energy of the noninteracting cooperons and bogolons. After summing up over κ, from (25) we obtain that the total quasi-particle energy in the BCS approximation can be rewritten in the following form

$$\tilde{E} = \Sigma 2\xi_k u_k^2 f_k + \Sigma 2\xi_k v_k^2 (1-f_k) - \Sigma \frac{\Lambda_k^2 (1-2f_k)}{2\lambda_k}. \quad (34)$$

The total quasi-particle energy of the noninteracting cooperons and bogolons is determined by the expression

$$E = \Sigma 2\lambda_k u_k^2 f_k + \Sigma 2\xi_k v_k^2 (1-f_k). \quad (35)$$

Substituting $2\lambda_k u_k^2$ with $2\xi_k u_k^2 + \frac{\Lambda_k^2}{\lambda_k}$ in (35) we get

$$E = \Sigma 2\xi_k u_k^2 f_k + \Sigma 2\xi_k v_k^2 (1-f_k) + \Sigma \frac{\Lambda_k^2 f_k}{\lambda_k}. \quad (36)$$

The difference between $\tilde{E}$ and E, determined by the equation

$$\tilde{E} - E = -\Sigma \frac{\Lambda_k^2}{2\lambda_k} = -0.5 N(\mu) \Lambda^2 \left(\frac{\hbar\omega_D}{\mu}\right) \quad (37)$$

coincides with the change of the total energy of the zero oscillations of the lattice /30, p. 161/. As the term, describing the phonon energy is dropped out in the reduced BCS Hamiltonian (6), the change of the energy spectrum of the zero oscillations cannot be taken into account. In such a way, there is a full coincidence of both theories, the BCS theory and the two-fluid theory, as far as the energy description is concerned. Replacing the sum by an integral in (35) we get that the change of the total quasi-particle energy in the transition of the metal from a normal to a superconducting

state at $T = 0°K$ is determinated by the following formula

$$E_S - E_N = \frac{N(\mu)}{\sqrt{\mu}} \left[2 \int_0^{\hbar\omega_D} \sqrt{\mu - \xi}\, \xi\, d\xi - \right.$$

$$\left. - \int_0^{\hbar\omega_D} \sqrt{\mu - \xi} \left(1 + \frac{\xi}{\sqrt{\xi^2 + \Delta^2}}\right) \sqrt{\xi^2 + \Delta^2}\, d\xi + \int_0^{\hbar\omega_D} \sqrt{\mu + \xi} \left(1 - \frac{\xi}{\sqrt{\xi^2 + \Delta^2}}\right) \xi\, d\xi. \right]$$

(38)

As the domain of integration in (38) is very small compared to the value of the chemical potential, then the following inequality $\Delta \ll \hbar\omega_D$ takes place and therefore with a sufficient precision one may ignore the dependence of the density of states on the quasi-particle energy ξ. In such an approximation of constant density of states from (38) we obtain the familiar result

$$E_S - E_N = -N(\mu)\Delta^2/2. \qquad (39)$$

12. Physical Description of the Correlation between CQEs Forming the Cooper Pair

The Bogolubov-Valatin transformation shows, that the new quasi-particles forming the cooper pair are fermions too. The assertion for binding cooper pairs has no physical sence, because the virtual exchange of phonons ensures a correlation between both cooperons of the pair only. The Bogolubov-

Valatin transformation separates the bogolons with the cooperons. The Cooper's result ought to be understood in such a sence, that the maximum contribution in the decrease of the total quasi-particle energy of CQEs is brought by the virtual exchange of phonons between CQEs pairs, having opposite spins and equal, but opposite wave vectors. Only for virtual exchange of phonons a simultaneous transition of two particles of the pair from one state to another is necessary.

When phonons are exchanged between CQEs both the energy and the quasi-momentum conservation laws are observed. But in scattering by exchange of a virtual phonon the energy conservation law is observed only for the transition from the initial to the final state. It is not necessary that the energy conservation law should hold for the transition from the initial to the intermediate state (i.e., to a state in which the first CQEs has emitted a phonon already, but the other CQE has not absorbed this phonon jet. This abeyance of the energy conservation law takes place because of the uncertainty relation between energy and time $\Delta t . \Delta E \geq \hbar$. If the life-time of the intermediate state τ is very small, then the uncertainty in the energy is large. In the process of a virtual emission and absorption of a phonon the energy is conserved with an accuracy of the order $\Delta E = \hbar/\tau$. Such virtual processes are possible only when a second CQE exists, which is capable to absorb the emitted phonon from the first CQE at once.

The virtual exchange of phonons provides a strong correlation between CQEs, since

during this exchange they cannot take part in other scattering processes. This is the reason the life-time of the intermediate state $\tau = \hbar/\Delta$ to be called mean free part time of CQEs, and the corresponding distance passed by them during the process $\ell = \hbar v_F/\Delta$ - the mean free path. Naturally, the distance, at which the CQEs are at the end of the process (in units 2π) may be called the coherence length ($\xi_0 = \ell/\pi$) /30, p.170/. On the other hand, the coherence length coincides with the uncertainty in the instantaneous position of the CQE pair in the process of a virtual exchange of a phonon ($\delta x \frac{\Lambda}{v_F} \sim \hbar$) /31/. It follows from the above discussion, that the regidity of the correlation is caused by the energy gap between the spectra of bogolons and cooperons. On the one hand, the energy gap limits the life-time of the intermediate state, but on the other, it serves as an obstacle to the small change of the total quasi-momentum.

Naturally, when other independent channels of scattering with mean free path ℓ_i exist, the mean free path of CQEs in s superconductor is determined by the familiar formula /24, p.59/

$$\frac{1}{\xi} = \frac{1}{\xi_0} + \sum \frac{1}{\ell_i} . \qquad (40)$$

In order to determine the CQEs' behaviour in external field it is necessary to take into account the probability of their free from scattering run, since during the scattering process the CQEs lose the whole information about the action of the external

field on them. In a superconductor the change in the behaviour of a CQE between the virtual exchanges of two successive phonons corresponds to the scattering of CQEs in a normal metal, because during the virtual exchange of a phonon the exchanging CQEs cannot take part in other processes. In this way the connection between the depth of penetration λ of external magnetic field in a superconductor and the mean free path of the CQEs become clear. Also the nonlocal Pippard relation between superconducting current $J_S(r)$ and vector-potential $A(r)$ can be obtained from Reuter-Sonfheimer's formula with the help of the mean free path /32/.

So, the experiments demonstrate, that the coherence length is a large, but finite quantity. Therefore, it is reasonable to assume that the CQEs, which participate in the virtual exchange of phonons, occupy states in different energy branches, separated by a finite energy gap. Actually, in case that the CQEs, participating in virtual exchange of phonons, occupy states, lying on the same energy branch, then the uncertainty of their energy would be insignificant, and the life-time of the intermediate state and the coherence length - infinitely large. From what has been said it follows, that the energy gap between the energy spectra of bogolons and cooperons plays a stabilizing role, so that the life-time of the intermediate state and the coherence length are decreased by its increase. One can see from here that the boundaries of the superconducting sample and the impurities,

included within it, restrict the possibility of the CQEs to pass without scattering the coherence distance, which is necessary for virtual exchange of a phonon.

13. Ultra-Sound and Super-High-Frequency Light Absorption

By means of the above arguments one can easily explain the absorption of ultra-sound [33] and super-high-frequency light [34]. These processes can be considered qualitativly without consideration of the temperature influence. Therefore we may confine with the consideration of these processes at $T = 0°K$ only, when the occupation number function f_k has a step function form with a jump at the Fermi level. Then all cooperon states u_k^2 below the Fermi level and all bogolon states v_k^2 above the Fermi level are occupied. As far as we know, the absorption of ultra-sound of SHF-light stimulates the transition of CQEs from superconducting to normal states. Although only the cooperons are correlated, however for the virtual exchange of phonons the simultaneous transition of two quasi-particles is necessary. Hence it follows, that for absorption of ultra-sound of electromagnetic radiation minimum two quasi-particles may transfer from one energy branch to another one [13, p. 120], [21, p. 320]. Therefore the absorption of ultra-sound or SHF-light is possibile only if the energy of the absorbed quantum is equal or larger than the doublet gap energy. This energy is necessary for moving simultaneously at least two quasi-particles

from the cooperon energy branch to the bogolon one, or vice versa. Therefore the lower limit of the absorption frequencies coincides with the double thickness of the energy gap.

14. Tunneling of Bogolons in M-I-S and S_1-I-S_2 Systems

By means of the above arguments one can explain the tunnelling of CQEs between a metal and a superconductor or between two superconductors, separated by a dielectric layer. As the mechanism of the CQEs tunnelling in the two-fluid model with semiconducting energy spectrum of the excitations has been described long ago /7,p.183/, /19,p.63/,/35,p.89/, it remains only to connect it with the energy spectra E_k and ϵ_k and the distributions of cooperons $u_k^2 f_k$ and bogolons $v_k^2 (1 - f_k)$ correspondingly. Since at $T = 0°K$ the empty bogolon states below the Fermi level and the empty cooperon states above the Fermi level do not take part in tunnelling processes, we shall neglect them inwhat follows and for the sake of simplicity the empty states are not shown on the figures.

In the normal state of the metal, the abrupt boundary of the energy distribution of CQEs coincides with the Fermi level. In the superconducting state the virtual exchange of phonons between CQEs makes the upper bound of their distribution in K-space sloping (fig. 3) and a gap appears in their energy distribution (fig. 4). The thermodynamic equilibrium between two normal metals, separated by a very thin dielectric layer,

occurs when the corresponding Fermi levels coincide. At T=0°K the tunnelling of cooperon pair through a very thin dielectric layer between two superconductors (Josephson effect [36]) when there is no potential difference between them shows, that the thermodynamic equilibrium occurs when the boundaries of the corresponding energy distributions of the cooperons coincide (as shown on fig. 5), because only they freely take part in the transport phenomena.

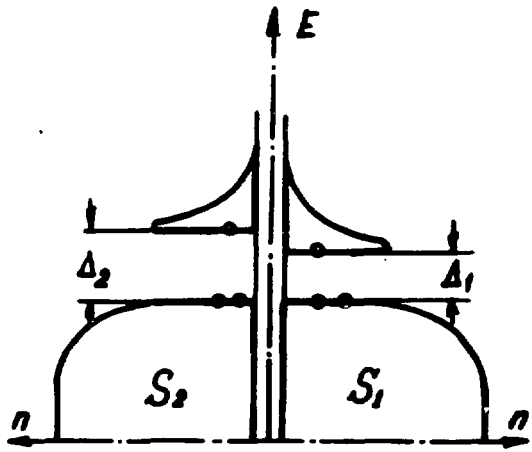

Fig. 5. Energy distribution of cooperons and bogolons in S_1-I-S_2 system at T = 0°K.

Although the bogolons have the energy spectrum ϵ_k of normal CQEs, they cannot take part in the transport phenomena, because they occupy all admissible states above the Fermi level. If the thickness of the dielectric layer between two superconductors increases,

the tunnelling of cooperon pairs decreases and there exists a critical thickness, above which the Josephson effect cannot be observed. If the thickness of the dielectric layer is greater than the critical one, only the transition of bogolons across the layer is possible. The absence of a tunnel current in this case at $T=0°K$ proves our physical supposition, that bogolons, situated above the Fermi level cannot take part in the transport phenomena at zero temperature. At $T=0°K$ all bogolons are forced to take part in the continuous virtual exchange of phonons, in order to provide the cooperons energy spectrum E_k. If on the one side of the dielectric layer there is a normal metal, while on the other side of it there is a superconductor [37], then the thermodynamic equilibrium takes place when the chemical potential of the CVE in the normal metal and the upper boundary of the cooperon energy distribution in the superconductor are equal (fig. 6). A tunnel current of normal CQEs appears at a potential difference $V = \Delta/e$ between the normal metal and the superconductor. This potential difference is necessary for the raising of CQEs in the normal metal up to the lower boundary of bogolon energy distribution in the superconductor, in order to ensure the free transition of CQEs from normal state of the metal across the dielectric layer to the bogolon states of the superconductor (fig. 7); or else in order to provide a simultaneous transition of two cooperons into bogolons, so that one of them tunnelling in the normal metal transfers part of its energy to the others. This energy in ne-

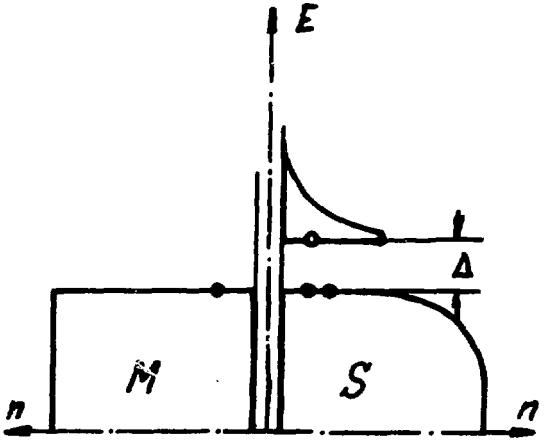

Fig.6.Energy distribution of cooperons,bogolons and normal CQEs in M-I-S system at T = 0°K.

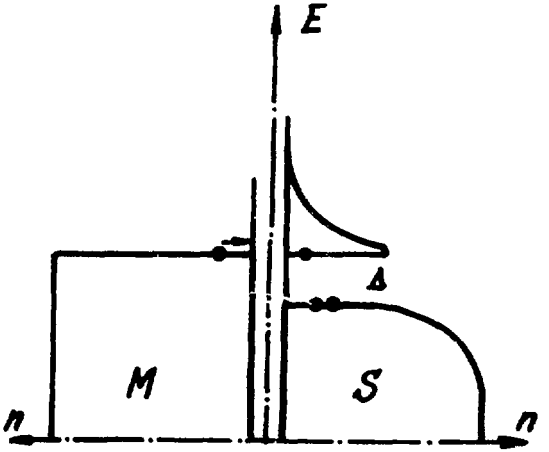

Fig.7.Energy distribution of cooperons,bogolons and normal CQEs in M-I-S system at V = Δ/e.

cessary for the donatory CQE so that it can pass to a bogolon state (fig. 8). If on both sides of the dielectric layer two superconductors are situated, then the tunnel current of bogolons appears at a potential difference of $V=(\Delta_1+\Delta_2)/e$. The letter is initially necessary to raise the upper boundary of the cooperon energy distribution of the one superconductor to the lower boundary of the bogolon energy distribution of the other one (fig. 9 and fig.10); and afterwards also for the transition of one cooperon of the same superconductor to a bogolon of the other one, giving part of its energy to the corresponding cooperons. This energy is necessary for its transition to a bogolon state in the superconductor with the raised energy distribution (fig. 11 and fig. 12). Similar mechanism of tunnel transition of normal CQEs is already used for description of the tunnel current in the semiconducting model of the energy spectrum of the excitations. But in this work cooperons have a semiconducting energy spectrum (15), while bogolons have the normal energy spectrum (1). Therefore it is necessary to use the new density of states of cooperons and bogolons, corresponding to new energy spectra, so that the experimental data correspond perfectly to the given theory.

15. Conclusion

Thus the physical picture of the metal transition from a normal to a superconducting state had been described in the above suggested two-fluid model of the supercon-

ducting quasi-particles. The choice of the Bogolubov-Valatin transformation parameters by means of bogolons and cooperons divides the CQEs only, but cannot influence upon their energy spectra. In conformity with the accepted terminology in BCS theory it is reasonable to ascribe to the correlated quasi-particles the name cooperons, and to the uncorrelated - bogolons. Cooper's idea for bound pairs, forming Bose-condensate, is unwarrantable, in spite of its usefulness for description of the superconductivity, because the coherence length is inconceivable larger in comparison with the average distance between them. The energy gap between the energy spectra of cooperons E_k and bogolons ϵ_k, together with the strong correlation between the CQEs, preventing the small

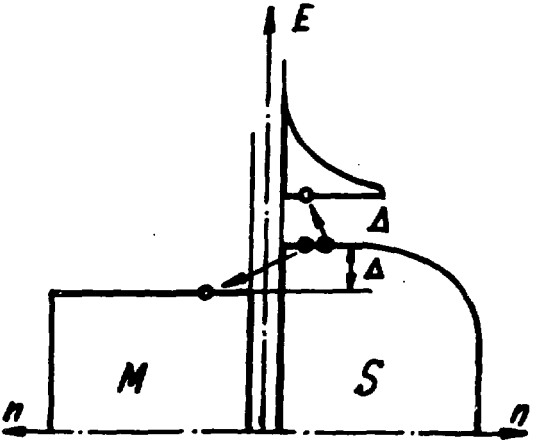

Fig. 8. Energy distribution of cooperons, bogolons and normal CQEs in M-I-S system at $T = 0°K$ and $V = -\Delta/e$.

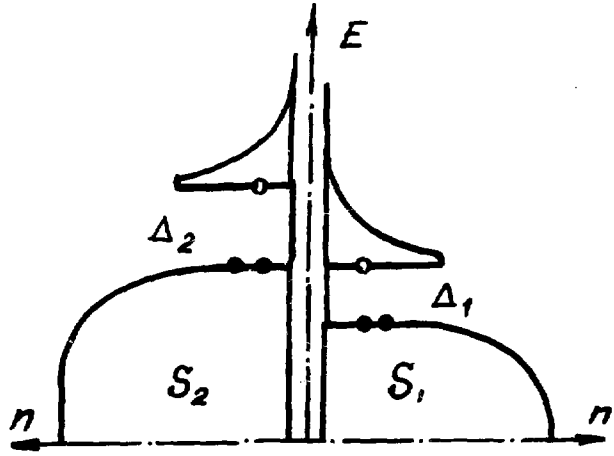

Fig.9. Energy distribution of cooperons and bogolons in S_1-I-S_2 system at T=0°K and $V=\Delta_1/e$:

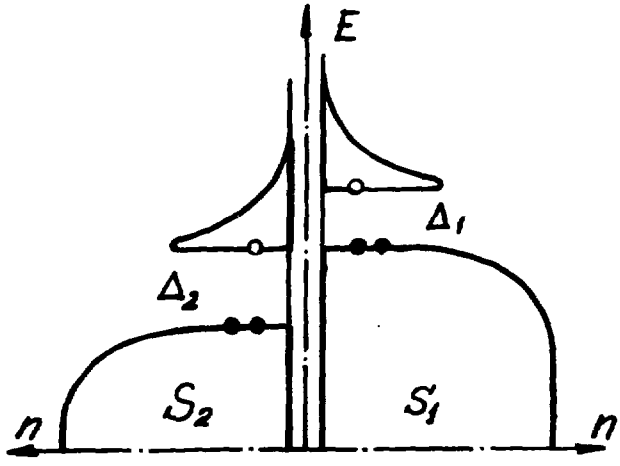

Fig.10. Energy distribution of cooperons and bogolons in S_1-I-S_2 system at T=0°K and $V=-\Delta_2/e$.

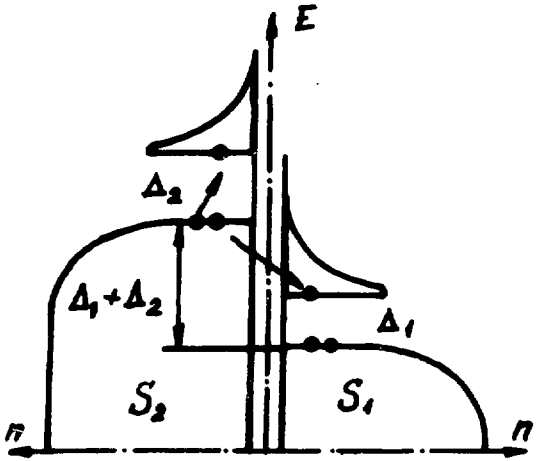

Fig.11. Energy distribution of cooperons and bogolons in S_1-I-S_2 system at T=0°K and $V=(\Delta_1+\Delta_2)/e$.

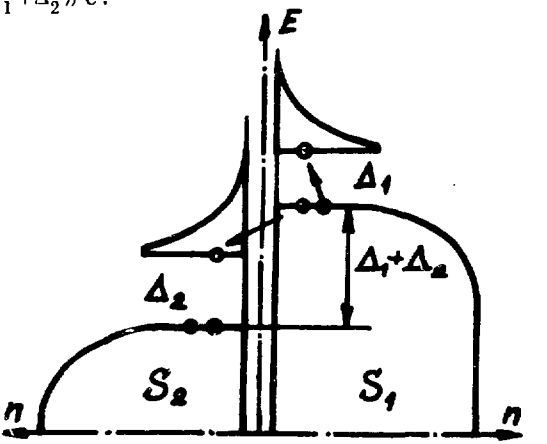

Fig.12. Energy distribution of cooperons and bogolons in S_1-I-S_2 system at T=0°K and $V=-(\Delta_1+\Delta_2)/e$.

change of the total quasi-momentum ensure infinite conductivity of a direct current in the superconductors /7, p.172/, /38/.

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