The application of a strong magnetic field creates the following advantages:

1) For an isolated exciton, when $\hbar\Omega \gg \epsilon_0$, where $\Omega = eH/m^*c$ is the cyclotron frequency, the binding energy is of the order of $\epsilon_0 \ln^2(\hbar\Omega/\epsilon_0)$. In the Brandt-Chudinov experiments $\hbar\Omega/\epsilon_0 \sim 10^4$, so that the binding energy became of the order of several degrees.

2) In the metallic limit, i.e., for $E_g < 0$, it is sufficient for the strong magnetic field to "one-dimensionalize" the motion of the electrons and holes. In consequence, the deleterious effect of the spectral anisotropy is completely eliminated, and the formation of an exciton dielectric is possible at any carrier density, i.e., at any $E_g < 0$.

The theoretical investigation in the high-density limit has revealed many different possible cases, depending upon the direction of the magnetic field and the sign of the effective interaction between the carriers. Of extremely great help in this classification and analysis was the method developed by S. A. Brazovskiĭ for taking into account the transverse motion of the carriers, a motion which is described by zero-point oscillator functions.

As a result of the theoretical analysis, the following general conclusions were drawn:

a) Pairing of carriers of the same sign, i.e., superconductivity, is impossible.

b) If the effective interaction has the same sign as the Coulomb interaction, then pairing of electrons with holes, or of a quasiparticle of the electron type with a quasiparticle of the hole type from different electronic groups is possible if all these groups are not identical. Pairing occurs between those two carrier groups that interact most strongly. The rest remain free.

c) If the sign of the interaction is determined by the phonons, i.e., is opposite to that of the Coulomb interaction, then the pairing of a quasiparticle of the electron type with a quasiparticle of the hole type to form one electronic group is possible, but only if there are a few symmetric (with respect to the direction of the magnetic field) electronic groups. The holes and those electrons which do not pertain to symmetric groups remain free.

d) If the sign of the interaction is Coulombic and there are several symmetric electronic groups, then in the event of pairing of electrons from these groups with holes or with a nonsymmetric electronic group, even if all the electrons of the symmetric groups participate in the pairing, the physical properties of the system are such as if only one of the symmetric groups participates and the rest remain free.

e) If the direction of the field is nonsymmetric and all the electron groups are not identical, then the following sequence of transitions is possible: 1) the pairing of two groups, 2) the pairing of two of the remaining groups, etc.

The mathematical apparatus of the theory is similar to the theory of superconductivity. The physical properties of the material in the presence of pairing are determined by the fact that as the temperature is lowered from the critical temperature (the second-order phase transition point in a metal) $T_c \sim (p_0^2/m^*)\exp(-\kappa\hbar v/e^2)$ to zero, a portion of the carriers ceases to participate, mainly according to the law $e^{-\Delta/T}$, where $\Delta \sim T_c$. This affects the thermal and electrical conductivities, the electronic heat capacity, etc. In particular, the electrical conductivity takes the form

$$\sigma(T \ll T_c) = \sigma(0) + aTe^{-\Delta/T}.$$

The ratio $\sigma(0)/\sigma(T_c)$ depends on the specified case. If the effective masses of the electrons and holes satisfy the inequality $m_e \ll m_h$, then the conductivity is determined mainly by the electrons. If a fraction α of all the electrons remains after the pairing, then $\sigma(0)/\sigma(T_c) = \alpha$. If, however, all the electrons pair off, then $\sigma(0)/\sigma(T_c) \sim (m_e/m_h)^2 \ll 1$. The transition temperature T_c and the energy "gap" Δ decreases upon introduction of impurities. When the impurity concentration is higher than a certain critical value, such that the reciprocal collision time becomes equal to $\hbar/\tau_c \sim T_{c0}$ (T_{c0} is the critical temperature of the pure substance), no exciton dielectric is formed. At smaller concentrations (such that $0.91/\tau_c < 1/\tau < 1/\tau_c$) we get $T_c \neq 0$ but $\Delta = 0$, i.e., a phase is formed, but it is a "gapless" phase. In the presence of pairing, the appearance of a new small-amplitude periodicity of the potential in the crystal should be observed.

These theoretical predictions are fully confirmed by the experiments of N. B. Brandt and S. M. Chudinov (the highest critical temperature reached is 7°K). Furthermore, a second-order phase transition is experimentally observed when $E_g > 0$, i.e., from the dielectric phase. The theory of this phenomenon has not as yet been constructed.

I. B. Levinson and É. I. Rashba. Bound States of Electrons and Excitons with Optical Phonons in Semiconductors

A variety of the properties of solids is determined by the dispersion laws of quasiparticles and the nature of the interaction between them. Therefore, the appearance of quasiparticle bound states changes essentially the various properties of crystals, especially the optical properties. A well-known example is the Mott exciton (an electron-hole bound state).

Bound states appear below the disintegration threshold. The situation near the threshold of the decay in which an optical phonon is emitted is shown in Fig. 1. Above the threshold, $\epsilon = \hbar\omega_0$, the decay is possible, and there is no spectrum in this region. Therefore, when the "bare energy" $\epsilon_0(p)$ approaches the threshold, specific distinctive features appear in the intrinsic spectrum $\epsilon(p)$[1]. The approach of the energy ϵ_0 to the threshold $\hbar\omega_0$ can also be realized when the external parameter controlling the spectrum (magnetic field, pressure) is varied. In this case the threshold

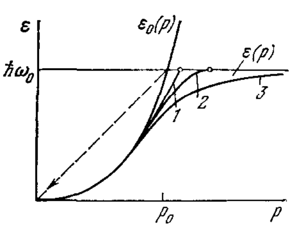

FIG. 1

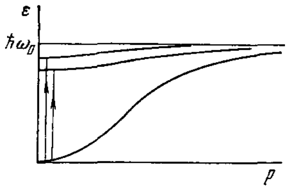

FIG. 2

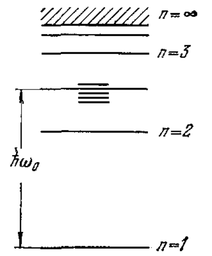

FIG. 3

situation corresponds to a resonance situation: the distance between two electronic and exciton levels coincides with $\hbar\omega_0$. It is precisely by such means that the threshold phenomena were experimentally observed[2,3]. Henceforth we shall take p to mean any parameter that controls the spectrum. Curves 1—3 show three typical ways in which the spectrum can behave: intersection with the threshold, tangency, and asymptotic approach. In the last two cases it is customary to speak of pinning[2,4-8].

The study of the interaction with optical phonons is interesting, in that one can follow the variation of the fractional phonon participation in the formation of a quantum state as p is varied. A curve of the type 3 arises in the magnetopolaron problem (p is a magnetic field[5,6] or a component of the momentum along H [7]) and in the impurity-center problem[4]. When p is appreciably larger than p_0 we have a bound state of a phonon of momentum $q \approx p$ and an electron of momentum close to zero. In other words, the phonon has almost the entire energy and momentum of the excitation, and the electron determines only the charge. It is clear that when the interaction between the particles is weak (the coupling constant $\alpha \ll 1$) it is impossible to imagine another excitation with energy close to $\hbar\omega_0$ and not depending on p. When $p \approx p_0$ the fraction of the phonon falls to $\frac{1}{2}$ (a hybrid state), while when p is appreciably less than p_0 the fraction falls to $\alpha \ll 1$. The character of the states along a curve of the type 2 changes in similar fashion, the bound states occurring near the end point. Such a curve arises, for example, in the magnetoexciton problem (p is the magnetic field[8]). On the other hand, for a curve of the type 1 the phonon participation is everywhere small and there are no bound states. It follows from the foregoing that in the majority of cases the states below the threshold turn out, as usual, to be the bound states of those particles which exist as free particles above the threshold.

The most significant theoretical result of the study of the spectrum near the threshold is the proof of the existence in this region of secondary spectral branches which are bound states over the entire spectral range. Such branches for a magnetopolaron are shown in Fig. 2; they form a sequence, crowding toward the threshold and lying at a distance $\sim\alpha^2\hbar\omega_0$ from it[9]. The experimental search for such electron and phonon bound states is an interesting problem, quite practicable, as estimates show, in CdTe, for example. Bound states should appear in the absorption spectrum of free carriers in the form of discrete lines (the transitions indicated by arrows in Fig. 2). Similar branches exist below the threshold for three-particle hole, electron, and phonon) production in a strong magnetic field[10]. The role of the magnetic field in these cases consists in creating a quasi-one-dimensionality—a high density of states at the bottom, of the band, where the electron falls after emitting a phonon (the arrow in Fig. 1). The decay probability, i.e., the effective force of interaction, increases as a result, and this favors the formation of a bound state.

Another factor that favors the formation of quasiparticle bound states is the suppression of the recoil kinetic energy for large quasiparticle masses. Therefore, for strong coupling ($\alpha \gg 1$), when the polaron mass is proportional to α^4 and is very large, and a sequence of polaron and phonon bound states arises in the absence of a magnetic field[11]; it is precisely in this problem that the secondary branches were first obtained. The recoil is completely suppressed if the electron is localized on an impurity center. Therefore, even in a weak coupling there arises an infinite system of levels (Fig. 3) describing the electron and phonon bound states[12]. In contrast to the magnetopolaron problem, a finite number of such levels lies above the threshold (corresponding to the number of electronic levels below the threshold). It is possible that bound states of this sort have already been experimentally observed[13].

It must be emphasized that bound states with optical phonons arise in a problem in which the number of particles is not conserved and, therefore, no simple model of the interaction in configuration space exists. In this connection the secondary branches can be found only in the study of the integral equation for the scattering amplitude or the wave function.

Notice also that the appearance of secondary branches is in no way connected with a resonance situation.

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