DESTRUCTION OF LONG-RANGE ORDER IN ONE-DIMENSIONAL AND TWO-DIMENSIONAL SYSTEMS POSsessing A CONTINUOUS SYMMETRY GROUP. II. QUANTUM SYSTEMS

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The asymptotic behavior of the correlations in the low-temperature phase are found for the following two-dimensional quantum systems: a two-dimensional lattice of plane rotators, two-dimensional magnetic moments on a plane of magnetization, a two-dimensional Bose liquid (Sec. 1), and a two-dimensional Ising ferromagnet (Sec. 2). The change of the asymptotic behavior with increasing temperature is investigated for the systems considered in Sec. 1. The nature of the low-temperature phase for a continuous symmetry group is also considered in Sec. 1. The nature of the low-temperature phase for a continuous symmetry group is also considered in Sec. 1.

I. In the previous article (1) by the author, the correlation functions were determined at low temperatures in two-dimensional systems of classical spins and to other two-dimensional systems possessing a continuous symmetry group. In the present article similar results are derived for two-dimensional quantum systems in Sec. 1, the correlations are investigated in a two-dimensional lattice of plane rotators, in a two-dimensional Bose liquid, and in planar magnetic systems in Sec. 2.

1. The nature of the temperature-dependent corrections to the found expressions is briefly described in Sec. 1, and the question of the nature of the low-temperature phase, which was treated incorrectly in article (1), is touched upon: The characteristic of the low-temperature phase is an infinite susceptibility, as was erroneously asserted in article (1), but rather the superfluidity of a two-dimensional Bose liquid and the corresponding property of two-dimensional magnetic substances. The author hopes to return to this question, too.

1. ONE-DIMENSIONAL LATTICE OF PLANAR ROTATORS AND SYSTEMS EQUIVALENT TO IT (TWO-DIMENSIONAL BOSE LIQUID AND PLANAR MAGNETIC SUBSTANCES)

Lattice of Planar Rotators

This system corresponds to a two-dimensional lattice with lattice constant a (the notation for the lattice sites is the same as in 1), a planar rotator described by the real variables \( \phi \) (angle) and \( \theta = \theta/\theta_{\text{B}} \) (angular momentum) is associated with each site \( r \), and the function of the system must be periodic with respect to each of the variables \( \phi \) with period \( 2\pi \), and the eigenvalue \( \phi \) are integers. We take the Hamiltonian of the system in the form

\[
\mathcal{H} = \sum_{r} \mathcal{H}_{r}(\phi) = \frac{1}{2} \sum_{r} \left( \frac{\partial^{2}}{\partial \phi^{2}} - \frac{\partial^{2}}{\partial \theta^{2}} \right) \phi_{r}^{2} + \frac{1}{2} \sum_{r} \left( \frac{\partial^{2}}{\partial \phi^{2}} + \frac{\partial^{2}}{\partial \theta^{2}} \right) \phi_{r}^{2}
\]

(1)

where only nearest neighbors interact (\( r - r' = \pm a \)). The energy interaction is an even periodic function of the difference \( \phi_{r} - \phi_{r'} \), having a single minimum at \( \phi_{r} = \phi_{r'} = 0 \) in the interval \( \phi_{r} < \phi_{r'} < \pi \), that is, the interaction is ferromagnetic:

\[
\mathcal{H}_{r} = \sum_{r' \neq r} J_{r} \phi_{r} \phi_{r'} + \mathcal{H}_{r}(\theta)
\]

(2)

\[
J_{r} = 1 - \sum_{r' \neq r} c_{r}^{2} \geq 0,
\]

(3)

For the time being we shall forget about the periodicity conditions on the wave function, regarding it as a function of the variables \( \phi_{r} \), which vary over the entire interval \( -\pi < \phi_{r} < \pi \). The potential energy then has minima for configurations satisfying the condition

\[
\phi_{r} - \phi_{r'} = 2\pi \frac{n}{N}
\]

(4)

where \( n \) is an arbitrary integer, i.e., defined on neighboring pairs of sites \( r \) and \( r' \) such that \( n \neq n' \). Each such configuration corresponds to the minimum configuration (4). We assume that a minimum one can replace the Hamiltonian by its quadratic expansion with respect to small deviations from (4), which gives

\[
\mathcal{H} \approx \mathcal{H}_{r}(\phi) = \sum_{r} \left( \frac{1}{2} \phi_{r}^{2} + \frac{1}{2} \theta_{r}^{2} \right)
\]

\[
= \frac{1}{2} \sum_{r} \left( \phi_{r}^{2} + \theta_{r}^{2} \right)
\]

(5)

We have a situation which is similar to the one which is encountered in the well-known method of Heitler and London, viz., in different regions of configuration space the Hamiltonian separates differently into the major part and the perturbation (so-called asymptotic perturbation theory, see (2)). In the first approximation the eigenfunctions of the initial Hamiltonian will be linear combinations of the eigenfunctions of the Hamiltonians (5).

In order to find these linear combinations it is, in general, necessary to solve the secular equation, but allowance for the symmetry requirements simplifies this problem. In the present case the total wave function

\[
\psi_{r}(\phi, \theta) = \sum_{n} c_{n} \phi_{r}^{n}(\phi)
\]

denotes the quasi-momentum

\[
\mathbf{p} = \sum_{n} c_{n} \mathbf{p}_{n}
\]

(6)

for classes \( \phi \) different from \( \phi_{n} \) is the Hamiltonian of a system of "dissolved" harmonic oscillators and its energy spectrum is shifted with respect to (8) by an amount \( \Delta \mathcal{E}_{\phi} \), equal to the minimum of the inhomogeneous quadratic form in (7). By virtue of this, the taking of the class \( \phi_{n} \) into consideration allows corrections of the order of exp(\( -\Delta \mathcal{E}_{\phi} \)) and at low temperatures one can confine attention only to the contributions from the class \( \phi_{n} \). We shall show that the average energy of the system averages reduces to the similar problem for a system of harmonic oscillators having the spectrum (8).

In fact, let the averaging with a Gibbs weight function (9) be carried out with respect to the approximate eigenfunctions (7) of the class \( \phi_{n} \) and the averaged quantity in a periodic local function \( \hat{\psi}_{n}^{(p)}(\phi) \) of the angles \( \phi_{r} \) then each term of the summation (7) gives the same contribution to the average, and it coincides with the average of \( \mathcal{H} \) over \( \hat{\psi}_{n}^{(p)}(\phi) \) over the system of oscillators. In particular, for integer values of \( m_{n} \) one can write

\[
\langle \mathcal{H} \rangle = \sum_{n} \langle \mathcal{H} \rangle_{n} = \sum_{n} \langle \mathcal{H} \rangle_{n} - \sum_{n} \langle \mathcal{H} \rangle_{n} \psi_{n}^{(p)}(\phi)
\]

(9)

where \( \langle \mathcal{H} \rangle_{n} \) is the correlation function for the coordinates of the macroscopic oscillators. One can easily deduce Eq. (9) for a system of harmonic oscillators for arbitrary values of \( m_{n} \), but, according to what was said above, in the case of the averaged function periodic, its average with respect to the system of oscillators coincides (for \( T < T_{c} < 1 \)) with the Gibbs average over the initial system.

For \( \psi_{n}^{(p)}(\phi) \) the correlation function \( \langle \mathcal{H} \rangle_{n} \) depends on the method of ordering the macroscopic operators \( \hat{\psi}_{n}^{(p)}(\phi) \), but at large distances these differences are these of the order of \( 1/T_{c} \), where \( \psi_{n}^{(p)}(\phi) \) is defined as

\[
\psi_{n}(\phi, \theta) = \sum_{r} \hat{\psi}_{n}^{(p)}(\phi)
\]

(10)

that is, the operators \( \psi_{n}^{(p)}(\phi) \) at remote points asymptotically commute.

For \( n_{0} \gg \alpha \) a large number of appreciable system in the order of magnitude of the eigenvalues of \( \mathcal{H} \), according to the which due to the strong dependence of the averaged function periodic, its average with respect to the system of oscillators coincides (for \( T < T_{c} < 1 \)) with the Gibbs average over the initial system.

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first term in (11) is a logarithmically divergent constant (for \( R \rightarrow \infty \), where \( R \) is the size of the system).

Substituting \( \langle n(0) \rangle \) and taking into consideration that the first term from (11) gives a factor
\[ \exp[-\mathcal{U}(\mathcal{A}) / (kT)] \]
the right-hand side of (10), which is equal to \( \sum_{n=0}^{\infty} \left( \frac{m}{m+1} \right)^n \), to the limit \( R \rightarrow \infty \), we obtain the following result for \( u_0 \) such that
\[ a_0 \equiv \frac{\sum_{n=0}^{\infty} \left( \frac{m}{m+1} \right)^n}{(m+1)^n} \]
are integers, the Laughlin estimates (12) become
\[ \alpha_s(\infty) \approx 2a_0 \]
where \( a_0 \approx 1.19 \) for all \( \alpha_s > 1 \).

Correlations involving the participation of \( \rho_{ij}(\mathbf{r}) \) are of the same order, \( Q_{ij}(0) \), as the sums associated with the noncommutativity of \( w_{ij}(\mathbf{r}) \), which were shown away earlier, so that at large distances one can neglect them and assume
\[ \rho_{ij}(\mathbf{r}) \sim \rho_0(\mathbf{r}) \]
the asymptotic properties of the correlation function. For \( t \rightarrow \infty \), expression (13) is equivalent to formula (21) from the previous section, but with a different value of the cutoff parameter \( \sigma \).

Two-dimensional Bose Liquid

Now let us show that the long-range correlations in a two-dimensional Bose liquid can be calculated on the basis of an equivalent Hamiltonian isomorphic to (1). Let us assume that the Bose condensate only exists in the ground state \( \mathcal{H}_{\text{B}} = \mathcal{H}_0 \), and for \( T > 0 \) it is \( \mathcal{H}_0 \), so that the fluctuations are determined by the quantum state of the condensate, which may be calculated by the formalism of the Bose liquid model is equivalent to \( \mathcal{H}_0 \) where \( \mathcal{H}_0 \) is a projection operator on a subspace spanned by the states \( \mid v \rangle \). More precisely
\[ \mathcal{H}_0 = \frac{1}{2} \sum_{\mathbf{r}, \mathbf{r}', \mathbf{r}''} \rho_0(\mathbf{r}) \rho_0(\mathbf{r}') \rho_0(\mathbf{r}'') \]
where \( \rho_0(\mathbf{r}) \) is the operator multiplying the state \( \mid v \rangle \) by the number \( \rho_0(\mathbf{r}) \) and \( \mathcal{H}_0 \) is the sum of all terms arising from the linearization of the Hamiltonian corresponding to transitions from cell to cell. The sum over the number of particles \( \rho_0(\mathbf{r}) \) is left. It is obvious that these matrix elements will differ from zero only for transitions \( \mid v \rangle \rightarrow \mid \nu \rangle \) such that \( \nu \) is a neighboring cell. We shall, therefore, in the following, consider only the number of particles in the remaining cells does not change. Let us denote the value of such a matrix element by \( \rho_0(\mathbf{r}) \) and then consider the whole operators \( \mathcal{H}_0 \) in the form
\[ \mathcal{H}_0 = \sum_{\mathbf{r}} \left( \rho_0(\mathbf{r}) - \rho_0(\mathbf{r}) \right) \]

where \( \rho_0(\mathbf{r}) \) denotes the operator that changes the number of particles in the cell corresponding to \( \mathbf{r} \). Then we shall consider that the number of particles in the remaining cells does not change. Let us denote the value of such a matrix element by \( \rho_0(\mathbf{r}) \) and then consider the whole operators \( \mathcal{H}_0 \) in the form
\[ \mathcal{H}_0 = \sum_{\mathbf{r}} \left( \rho_0(\mathbf{r}) - \rho_0(\mathbf{r}) \right) \]

The Nature of the Correlations to (13) and the Nature of the Low-Temperature Phase

The following approximations were made in the derivation of the asymptotic expressions (13) for the correlation functions: In the first place, the contributions from the class \( \beta_j \) of quartic interactions were neglected, i.e., quantized vortices were not considered. In the second place, the correlations from the class \( \beta_j \) were calculated to the order \( g^3 = \mathcal{H}_0 \) and the contributions were calculated. For the low temperatures, the contributions are calculated to the order \( g^4 = \mathcal{H}_0 \) and the contributions from the quartic interactions are neglected. In the three terms of the fourth-order and higher terms in the expansion of the Hamiltonian in powers of \( g^4 \), the following contributions are made to the leading order: From the quartic terms, \( \mathcal{H}_0 \) is the quadratic form in the coordinates of the vortices, \( \beta_j \) is the quartic form in the coordinates of the vortices, and the contributions from the quartic terms are neglected. In the three terms of the fourth-order and higher order expansions in the Hamiltonian in powers of \( g^4 \), the following contributions are made to the leading order: From the quartic terms, \( \mathcal{H}_0 \) is the quadratic form in the coordinates of the vortices, \( \beta_j \) is the quartic form in the coordinates of the vortices, and the contributions from the quartic terms are neglected. In the three terms of the fourth-order and higher order expansions in the Hamiltonian in powers of \( g^4 \), the following contributions are made to the leading order: From the quartic terms, \( \mathcal{H}_0 \) is the quadratic form in the coordinates of the vortices, \( \beta_j \) is the quartic form in the coordinates of the vortices, and the contributions from the quartic terms are neglected.
In addition, it is evident in Eq. (21) that \( C(T) \rightarrow 0 \) as \( T \rightarrow \infty \).

The authors of \(^2\) justified the necessity of a phase transition by the fact that at low temperatures, the susceptibility \( \chi = \langle \delta m^2 / m^2 \rangle \) is infinite (the average moment \( m \) in the field \( \theta \) is given by formula (49), see below). The transition to the usual dependence \( m \sim h \) occurs at \( T \sim 2 \). Although the value of the \( T_c \) is unknown, by taking the interaction with the external field in the form \( h = c \cos \omega \) we obtain a similar transition when \( c \sim h \). The critical exponents \( \gamma, \beta, \), and \( \Delta \) are the same for the two-dimensional system of \( m \) spins which was solved in \(^5\) (for low temperatures).

First let us consider a single spin. As is well known, the components \( x, y, z \) of the vector \( \mathbf{S} \) are the generators of the representation of the SU(2) group. The matrix \( \mathbf{S}(\theta) \) is the matrix which corresponds in a given representation to the element \( \mathbf{U} \) of the rotation group. One can parameterize the elements of the SU(2) group \( \mathbf{U} \) by the Euler angles \( \alpha, \beta, \gamma \) which are universal covering group of the rotations, in the same way as the elements of the rotation group \( \mathbf{U} = \mathbf{E}(\alpha, \beta, \gamma) \) is the SU(2) group.

Let us consider a two-dimensional Heisenberg ferromagnet. Let \( \mathbf{S}^2 = (2 S_z + \sum \mathbf{S}^2) \) denote the Hamiltonian of the system; we shall utilize the identity (30), having taken as \( \mathbf{Q} \) the ground state of the \( \mathbf{E}(2) \) (see below) in the following form (which is equivalent to the usual formalism): For any operator \( \mathbf{F} = \mathbf{F}(\mathbf{Q}) \) acting on the SU(2) representation space of the system, and for any normalized vector \( \mathbf{v} \) of this space, the following identity is satisfied:

\[
\mathbf{F}(\mathbf{Q}) \mathbf{v} = \sum_{\mathbf{Q}} \mathbf{F}(\mathbf{Q}) \mathbf{v} = \mathbf{F}(\mathbf{Q}) \mathbf{v}.
\]

Now we shall consider a lattice in which a spin \( \mathbf{S} \) is associated with each lattice site \( \mathbf{r} \) (all the spins have the same magnitude, \( \mathbf{S}_r = \mathbf{S} \)). We associate its own SU(2) transformation \( \mathbf{S}(\mathbf{r}) \) with each site \( \mathbf{r} \) and the corresponding SU(2) transformation with the initial SU(2) transformation, which operates only on the states belonging to the spin \( \mathbf{S}_r \). Each lattice point \( \mathbf{r} \) is characterized by the triplet of functions \( x(\mathbf{r}), y(\mathbf{r}), z(\mathbf{r}) \), i.e., the triplet of functions \( \mathbf{S}(\mathbf{r}) = \mathbf{E}(\alpha, \beta, \gamma) \) (the product runs over all spins \( \mathbf{S}_r \) which transform the states of the entire system. This transformation can be regarded as the same as a transformation of operators whereby each operator \( \mathbf{F}(\mathbf{S}) \) is replaced by \( \mathbf{F}(\mathbf{S}) \mathbf{v} \), which acts on the space of states of the system, corresponds to a transformed operator given by

\[
\mathbf{F}(\mathbf{Q}) = \mathbf{F}(\mathbf{S}^2) = \mathbf{F}(\mathbf{S}_1, \mathbf{S}_2, \ldots) = \mathbf{F}(\mathbf{S}_1) \mathbf{F}(\mathbf{S}_2) \ldots.
\]

If the operator \( \mathbf{F} \) is expressed in terms of \( \mathbf{S}_r \), then the transformed operator \( \mathbf{F}(\mathbf{S}) \) is given by this expression by replacing each of the \( \mathbf{S}_r \) by \( \mathbf{S}_r \).

One can describe the properties of the "two-dimensional crystals" considered in \(^3\) in a similar fashion. The components of the deformation are treated as the independent variables which serve to describe the deformed crystals. Transverse waves exist in the system, so that the description describes a solid, notwithstanding the absence of long-range crystalline order.

\[\mathbf{F}(\mathbf{Q}) = \prod_{\mathbf{Q}} \mathbf{F}(\mathbf{Q}) \]

Substituting this expression into (32), we see that the quantum-mechanical average of the product of local operators \( \mathbf{A}_r \) pertaining to widely separated points \( \mathbf{r} \) coincides with the average over the equivalent classical system, in which the configurations are described by

\[\mathbf{E}(\mathbf{Q}) \mathbf{v} = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

The second term on the right-hand side of Eq. (37) can

\[\mathbf{E}(\mathbf{Q}) \mathbf{v} = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

In particular, here the components \( \mathbf{E}(\mathbf{Q}) \) of the spin are related to the classical quantities \( \mathbf{E}(\mathbf{Q}) \), where \( \mathbf{E} \) is a unit vector having the direction \( \mathbf{Q} \). Finally, in the spherical coordinate system:

\[\mathbf{E}(\mathbf{Q}) = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

In order to use the results of the investigation of the classical system \(^3\) (although the results are incorrect), it is only necessary to determine the expression for the equivalent energy \( \mathbf{E}(\mathbf{Q}) \). This connection is very important that the equivalent investigation \(^3\) was based only on the expression for the energy of slowly varying configurations, that is, configurations in which the difference between adjacent values is small:

\[\mathbf{E}(\mathbf{Q}) = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

One can find the expression for the equivalent energy of slowly varying configurations correct to within terms of second order in \( \mathbf{Q}^2 \) by using the corresponding expression for the transformed Hamiltonian:

\[\mathbf{E}(\mathbf{Q}) = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

Here \( \mathbf{Q} \) is the initial Hamiltonian \( \mathbf{H}_I \), \( \mathbf{H}_{II} \) and \( \mathbf{H}_{III} \) are the terms of first and second order in \( \mathbf{Q} \) and \( \mathbf{Q}^2 \), respectively.

If \( \mathbf{E}(\mathbf{Q}) \) is the initial Hamiltonian, then for the substitution \( \mathbf{S} \rightarrow \mathbf{S}^2 \), \( \mathbf{Q} \rightarrow \mathbf{Q} \), \( \mathbf{Q}^2 \rightarrow \mathbf{Q}^2 \), \( \mathbf{Q}^3 \rightarrow \mathbf{Q}^3 \) for each pair of neighboring spins \( \mathbf{r} \) and \( \mathbf{r}^2 \), the term of zero order in the expansion of \( \mathbf{E}(\mathbf{Q}) \) is then equal to \( \mathbf{E}(\mathbf{Q}) \), because the secular product does not change upon an identical rotation of both spins. Thus, the term of zero order in \( \mathbf{Q}^2 \) reduces to \( \mathbf{H}_{I} \) (which is simply a consequence of the equivalence of \( \mathbf{E}(\mathbf{Q}) \) with respect to the simultaneous rotation of all the spins in the system). One can obtain explicit expressions for the terms of first and second order in \( \mathbf{Q} \) and \( \mathbf{Q}^2 \) by using formula (29) and by calculating the expansion of \( \mathbf{E}(\mathbf{Q}) \) in powers of \( \mathbf{Q} \) based on the well-known expressions for \( \mathbf{E}(\mathbf{Q}) \) in terms of the Euler angles \( \alpha, \beta, \gamma \). We shall not present these calculations here, and we shall not present the final expressions for \( \mathbf{H}_{I} \) and \( \mathbf{H}_{II} \), limiting attention to a description of the method used to obtain these results.

In virtue of (36) one can put perturbation theory to calculate the values of the functional \( \mathbf{A} \) for slowly varying configurations, regarding the first term on the right-hand side of (36) as the major term and treating the terms \( \mathbf{H}_{II} \) and \( \mathbf{H}_{III} \) as perturbations. Thus, we arrive at an expression of \( \mathbf{E}(\mathbf{Q}) \) in powers of the differences \( \mathbf{Q} \), where the first-order term vanishes and the second-order term gives

\[\mathbf{E}(\mathbf{Q}) = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]

The second term on the right-hand side of Eq. (37) can

\[\mathbf{E}(\mathbf{Q}) = \mathbf{E}(\mathbf{Q}) \mathbf{v} \]
Electroacoustomagnetic Effect and Hall Effect in Semiconductors in a Strong Electric Field

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Large-scale properties of a semiconducting crystal in external electric and magnetic fields are considered for the case when the carrier drift velocity exceeds the phase velocity of the sound wave and phonon generation occurs. It is shown that the generated phonon current (besides the Lorentz force) acts on the electrons and which is, generally speaking, not a potential force. It is precisely the nonpotentiality of the electroacoustic force connected with the spatial anisotropy of the phonon emission diagram which leads to the appearance in the sample of an anomalous current component, and, as a consequence, formation of a magnetic moment (electroacoustomagnetic effect). The Hall effect in a strong electric field under phonon-generation conditions is considered for two limiting cases: for Hall-shorted and Hall-open samples. It is shown that in the case of a Hall-open sample the Hall constant decreases and reverses sign with increasing electric field (the absolute value of the Hall constant in this case may exceed a value in a field). The current-voltage characteristic is determined, and it is shown that under phonon-generation conditions the sign of current may reverse in a Hall-open sample. The case of a Hall-shorted sample with a Corbino disc geometry is considered and it is shown that in a strong magnetic field the current in the source circuit rises sharply. In weak magnetic fields, the current is saturated. The theoretical results are compared with available experimental data and good qualitative agreement is found. Kinematic effects connected with phonon generation are also mentioned.

Experim ental investigations of the Hall effect in semiconductors under a strong electric field have shown that it is accompanied by a number of new singularities and phenomena which are quite difficult to explain. Typical examples, for example, the monograph (10), most theoretical work on the Hall effect in a strong electric field initially begins with the premise that in a strong electric field the electron (hole) distribution function is significantly altered, so that the electric field can no longer be regarded as a small correction to the equilibrium distribution curve and depend on the electric field. By considering further some particular type of carrier scattering, with different dependences of the scattering time on the carrier energy, it is naturally possible to obtain different corrections to the Hall constant, necessitated by the action of the strong electric field. The galvanomagnetic properties and current-voltage characteristics of semiconductors in strong magnetic and electric fields were investigated in detail by Baranov and co-authors (11), who has shown that in a strong electric field, when the heating of the electrons is appreciable, the current-voltage characteristic of a semiconductor is determined, and it becomes possible to obtain negative differential resistance. Analogous phenomena were investigated earlier in plasma physics in the context of runaway electrons.

At the same time, there is another more mechanism whereby the Hall constant can be altered in a strong electric field; this mechanism is connected with generation of acoustic phonons by supersonic drift of the electrons. (10) Indeed, if a sufficiently strong electric field is applied to the sample, such that the average directional velocity of the electrons or holes exceeds the phase velocity of the acoustic wave, then, as is well known, acoustic phonons, sometimes also called acoustic noise, are produced in a crystal with relatively strong electron-phonon interaction. Owing to the acousticoelectric effect, the growing flux of acoustic phonons in the crystal affects the currents on the electrons, and an additional force produced by the generated phonons will act on the electrons in addition to the Lorentz force. If under these conditions the Hall effect and the Hall constant determined from it are significantly altered. Moreover, as will be shown below, a case is possible when the sign of the Hall constant is reversed (without a change in the type of carriers). Since the acousticoelectric force exerted on the electrons by the phonons generated in the crystal depends on the coordinates, it is natural that the experimentally measured Hall potential difference and Hall constant will also depend on the coordinates.

It should be noted that an investigation of the Hall effect in piezoelectric crystals p-Te recently carried out by Tsvetkov and Ioffe (12) has shown that in a strong electric field the Hall constant actually decreases and reverses sign, and that the region in which the indicated phenomena can be observed is close to the region of phonon generation. Apparently, in a number of other experiments, where a decrease of the Hall constant

(9) The influence of uneven distribution of the phonons in strong electric and magnetic fields was also investigated by Chernikov (7), who has shown that the mutual dragging of the electrons and phonons leads to a strong influence of the semiconductor. However, he considered only the kinematic motion of the electrons, and there was no phonon generation.

(10) The influence of the current of acoustic noise on various different conditions of the acoustic wave on different kinetic coefficients (particularly in the Hall effect in an acousticoelectric current) was considered by Gulyaev (7).