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EFFECT OF CHARGE ORDERING ON THE THERMOEMF OF LAYERED CRYSTALS IN A QUANTIZING MAGNETIC FIELD

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• *The paper deals with the effect of interlayer charge ordering on the thermoEMF of layered crystals in a quantizing magnetic field at helium temperatures for the case when the temperature gradient and the quantizing magnetic field are parallel to each other and perpendicular to the layers. Calculations are performed on the basis of a kinetic Boltzmann equation in the approximation of a stable relaxation time. It is shown that charge ordering in a quasi-classical magnetic field region leads to multiple sign inversion, the appearance of new thermoEMF oscillation frequencies and their amplitude rise at selected problem parameters by a factor of 5.9 – 13.5 as compared to the disordered state, which is caused by the conversion of crystal Fermi surface during phase transition due to charge ordering from a closed or transient surface into an open one. In strong magnetic fields, charge ordering leads to multiple switching of thermoEMF polarity with a change in a magnetic field. With selected problem parameters, maximum thermoEMF value increases by a factor of 7.6 – 18.4 as compared to the disordered state. However, in the ultra-quantum limit, the thermoEMF value is drastically reduced by the law* $\alpha \infty T^{-1}B^{-2}$*, which is due to the compression of crystal Fermi surface as a result of charge carrier condensation at the bottom of a single filled Landau subband with the number n = 0.*

Introduction. Current status of the problem

Nowadays, both radical improvement of conventional thermoelectric materials and intensive search for new ones is underway. Such conventional thermoelectric material as Bi_2Te_3 and other materials on its basis, as well as different aspects of their use in thermoelectric generators, coolers and measuring instruments are covered in comprehensive monographs [1, 2]. Apart from these materials, the investigations are concerned with metals, alloys, semiconductors [3, 4], fullerenes [5], composites [6], including biomorphic [7], etc. A promising thermoelectric material is graphene [8]. Its peculiar photothermoelectric effect, earlier considered to be purely photovoltaic, allows considering graphene as a material for high-efficiency solar elements.

Theory of thermoelectric properties of materials, including nanosystems, is being intensively developed [9, 10]. Thus, one of the first works dedicated, among other things, to theory of transverse thermoelectric coefficients of metals in a quantizing magnetic field, was published by A.M.Kosevich and V.V. Andreyev [11].

A large number of investigated materials, for instance, semiconductor systems of $A^{III}B^{VI}C^{VII}$ class, intercalated graphite compounds, synthetic metals based on organic compounds, graphene, etc. belong in their crystal structure to layered materials. At the same time, in the overwhelming majority of theoretical works dedicated to behaviour of such layered systems in quantizing magnetic fields the transverse galvanomagnetic effects are studied. The thermal conductivity of graphene in the presence of a magnetic field, including a quantizing magnetic field, its conductivity at direct and alternating current, the Hall quantum effect in it were considered in Ref. [12, 13]. In so doing, its Fermi surface was considered open, that is, such which occupies the entire one-dimensional Brillouin zone and, with

a periodic continuation, is a connected one, that is, represents a continuous corrugated cylinder. These works were mainly concerned with a research on charge carrier behaviour in the layer plane. Therefore, for the description of graphene band structure a concept of "massless neutrino" with a linear dependence of energy on the quasi-pulse in the layer plane proved to be efficient. In this case, the energy levels in a quantizing magnetic field were determined by the efficient Hamiltonian method.

In many cases, in layered crystals an interlayer charge ordering is observed. To such crystals belong fullerites [14], dichalcogenides of transient metals [15], synthetic metals based on organic compounds, etc. At the same time, thermoelectric properties of such materials, including those in charge-ordered state, have not been thoroughly studied so far. In this connection, the purpose of this work is theoretical study of the effect of charge ordering on thermoEMF of layered crystals in a quantizing magnetic field at low temperatures for the case when the temperature gradient and the quantizing magnetic field are parallel to each other and perpendicular to the layer plane.

Determination of magnetic field dependences of the electron gas chemical potential and charge order parameter

To calculate thermoEMF as a function of a quantizing magnetic field, one should know the field dependences of the electron gas chemical potential and the order parameter on the quantizing magnetic field. Note that in this paper we consider an interlayer charge ordering as alternation of layers with different electron density. Therefore, the surface density of electrons on the *i*-th layer is determined by the formula:

$$n_i = n_0 a\left[1+(-1)^i\,\delta\right]. \tag{1}$$

In this formula, n_0 is the average volumetric concentration of current carriers, a is the distance between translation equivalent layers, δ is the order parameter describing the nonuniformity of electrons in the layers, $\delta = 0$ corresponding to absence of ordering, and $\delta = 1$ – to full ordering.

The energy spectrum of charge carriers in a quantizing magnetic field, perpendicular to layers, is determined as follows [16]:

$$\varepsilon(n,x) = \mu^* B(2n+1) \pm \sqrt{W_0^2\delta^2 + \Delta^2\cos^2 x}\,, \tag{2}$$

In this formula, n is the Landau level number, $\mu^* = \mu_B m_0/m^*$, μ_B is the Bohr magneton, m^* is effective electron mass in the layer plane, B is magnetic field induction, Δ is initial half-width of mini-band that describes charge carrier motion perpendicular to the layers, W_0 is effective attraction between electrons that leads to charge ordering and is caused by a competence between the electron-phonon interaction and the Coulomb repulsion, $x = ak_z$, k_z is a quasi-pulse component in the direction perpendicular to layers.

With account of the energy spectrum (2), the chemical potential of charge carrier system ζ and the order parameter δ at low temperatures are found from the following system of equations [16]:

$$\begin{aligned}&\frac{1}{2\pi\zeta_0}\left[\int\limits_{\zeta+R\geq 0}(\zeta+R)\,dx + \int\limits_{\zeta-R\geq 0}(\zeta-R)\,dx\right] + \frac{kT}{\zeta_0}\sum_{l=1}^{\infty}\frac{(-1)^l}{\mathrm{sh}\left(\pi^2 lkT/\mu^* B\right)}\times\\ &\times\left[\int\limits_{\zeta+R\geq 0}\sin\left(\pi l\frac{\zeta+R}{\mu^* B}\right)dx + \int\limits_{\zeta-R\geq 0}\sin\left(\pi l\frac{\zeta-R}{\mu^* B}\right)dx\right] = 1,\end{aligned} \tag{3}$$

$$\frac{\delta W_0}{2\pi\zeta_0}\left[\int\limits_{\zeta+R\geq 0}(\zeta+R)R^{-1}dx-\int\limits_{\zeta-R\geq 0}(\zeta-R)R^{-1}dx\right]+\frac{\delta kTW_0}{\zeta_0}\sum_{l=1}^{\infty}\frac{(-1)^l}{\mathrm{sh}\left(\pi^2 lkT/\mu^* B\right)}\times$$
$$\times\left[\int\limits_{\zeta+R\geq 0}R^{-1}\sin\left(\pi l\frac{\zeta+R}{\mu^* B}\right)dx-\int\limits_{\zeta-R\geq 0}R^{-1}\sin\left(\pi l\frac{\zeta-R}{\mu^* B}\right)dx\right]=\delta. \qquad (4)$$

In these equations, $R=\sqrt{W_0\delta^2+\Delta^2\cos^2 x}$, $\zeta_0=n_0 ah^2/4\pi m^*$ is the Fermi energy of an ideal two-dimensional Fermi-gas at the absolute zero.

From equation (4) it follows that there always exists a trivial solution $\delta\equiv 0$ which corresponds to the absence of ordering. However, at $W_0/\zeta_0>1$ there also exists a nontrivial solution $\delta>0$ which corresponds to the ordered state. Condition $W_0/\zeta_0>1$ means that for the existence of the ordered state the electron-phonon interaction should not just exceed the Coulomb one, but this excess must be larger than the maximum kinetic energy of charge carriers in the layer plane with their given concentration. The reason for this lies in the increase of electron kinetic energy when passing to the ordered state.

The results of solving simultaneous equations (3) and (4) are shown in Figs. 1 and 2.

From Figure 1 it is seen that charge ordering reduces the chemical potential of charge carrier system, hence it is energetically favourable. Note that the chemical potential is counted from the middle of a pseudo-gap between mini-bands in the ordered state, which is matched by the middle of an output mini-band in the disordered state.

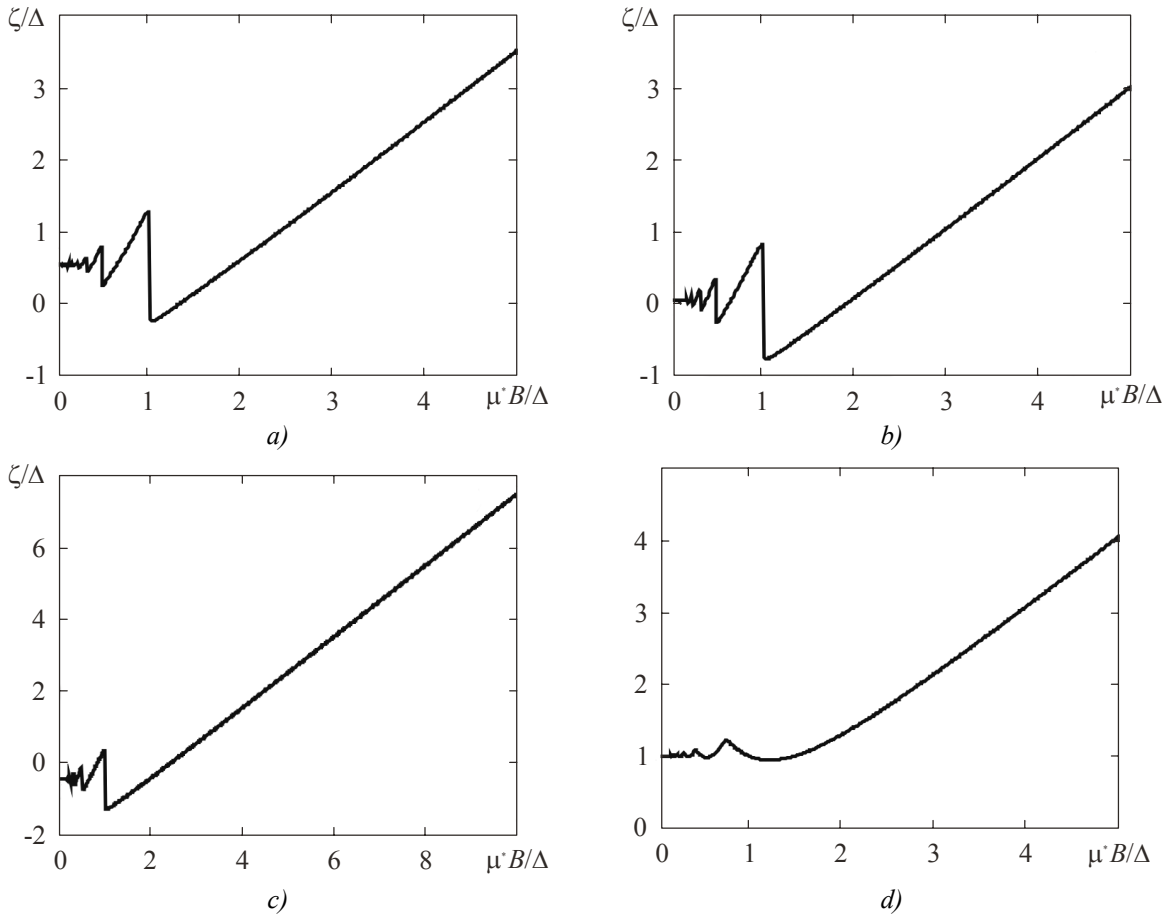

Fig. 1. Field dependence of the chemical potential of charge carrier gas at ζ_0/Δ = 1, kT/Δ = 0.03 and: W_0/ζ_0 = 1. 5 (a), 2 (b), 2.5 (c) and $\delta\equiv 0$ (d).

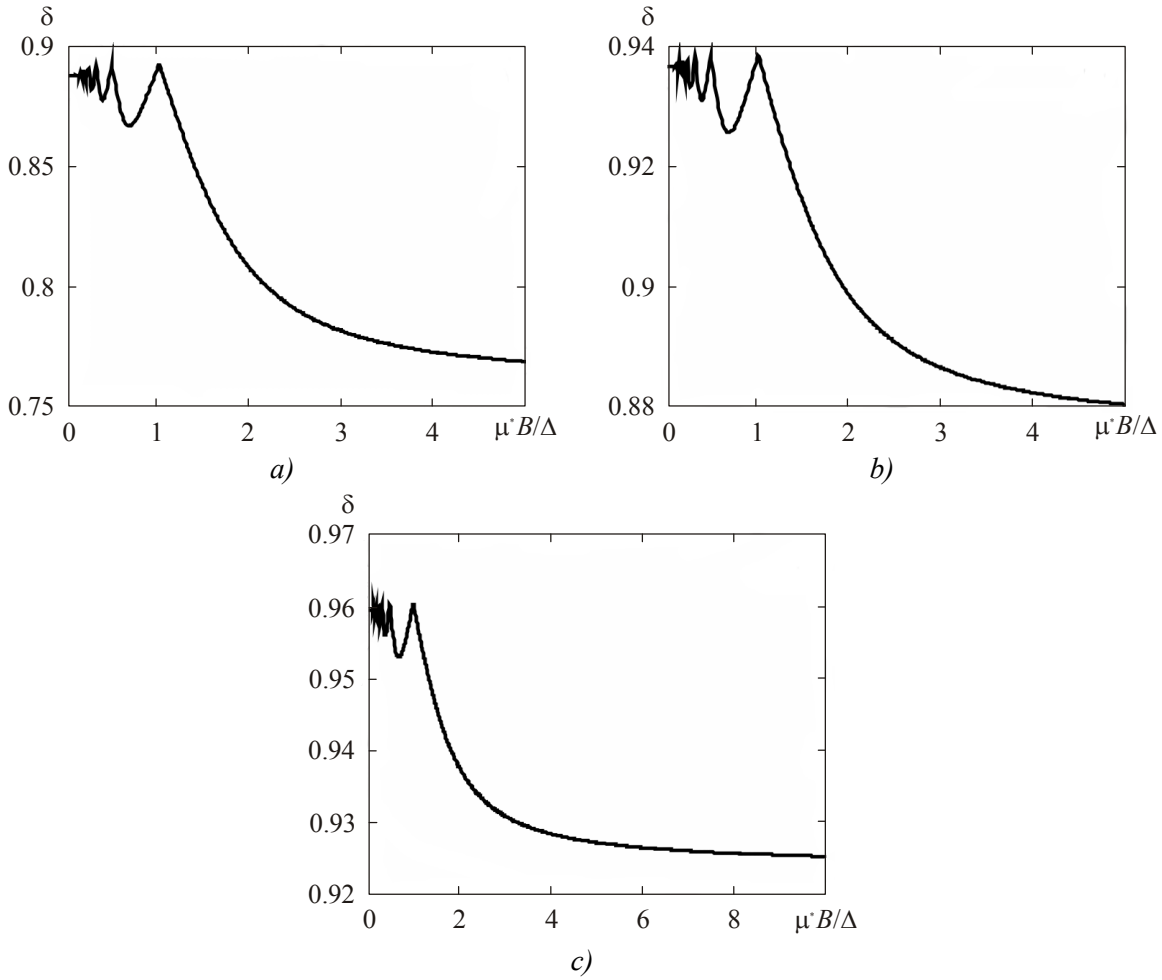

Fig. 2. Field dependence of the order parameter at $\zeta_0/\Delta = 1$, $kT/\Delta = 0.03$ and: $W_0/\zeta_0 = 1.5$ (a), 2 (b), 2.5 (c).

Transition into the charge-ordered state is a second-order phase transition and in the case when the Fermi surface of crystal in the disordered state is closed or transient, that is, when at $B = 0$ condition $-1 < \zeta/\Delta \le 1$ is met it is manifested as a topological transition from a closed surface into an open one. The topological transition is a transition whereby the Fermi surface changes connectivity, being converted from a closed surface into an open one, or vice versa. Recall that an open surface is such that occupies the entire single-dimensional (in k_z) Brillouin zone, appearing to be a connected surface with a periodic continuation. A closed surface is such that occupies only part of open Brillouin zone and with a periodic continuation appears to be composed of individual parts. From the figure it is also seen that, at first, the chemical potential in all the cases is an oscillating function of a magnetic field. At the same time, after the last oscillation maximum in the ordered state the chemical potential changes almost by hopping. Though, in fact, this hop is somewhat "smeared" due to the Dingle factor caused by current carrier scattering on the impurities and crystal lattice defects [17], but we will consider this factor to be close to unity. It can be done, if the ratio $l/a >> 1$, where l is mean free path at charge carrier scattering on the ionized impurities. The above hop testifies to a reversed topological transition from an open Fermi surface into a closed one. However, it should be noted that, strictly speaking, the Fermi surface of crystal at $B \neq 0$ is not defined, that is why it is referred to by convention, being imagined as a combination of hollow co-axial cylinders, the so-called magnetic tubes, "Landau tubes" whose axes are parallel to magnetic field direction. In this case the Fermi

surface is considered to be open, if at least one magnetic tube is found whose length is equal to the length of open Brillouin zone, and in the opposite case it is considered to be closed. After the last oscillation minimum, in all the cases, the magnetic field dependence of the chemical potential becomes almost linear. This dependence is due to the fact that at strong magnetic fields the chemical potential depends on the magnetic field as follows [16]:

$$\zeta = \mu^* B - \sqrt{W_0^2 \delta^2 + \Delta^2 \cos^2 \frac{\pi \zeta_0}{4\mu^* B}} \,. \tag{5}$$

In the derivation of formula (5) it was taken into account that the Fermi surface in the ultra-quantum limit, that is, under conditions of charge carrier condensation at the bottom of single lowest Landau sub-band, is composed of three parts. One of them is a continuous magnetic tube with symmetry plane $k_z = 0$ which narrows either side from this plane. The other two parts are halves of this tube and, starting from planes $k_z = \pm\pi/a$, they gradually narrow inward the open Brillouin zone. Thus, we really deal with a reversed transition from an open Fermi surface into a closed one in a strong magnetic field. As follows from formula (5), on the linear area of chemical potential field dependence parts of the Fermi surface corresponding to single filled Landau sub-band with the number $n = 0$, are gradually compressed in the direction of a magnetic field and this compression must be necessarily taken into account in the process of thermoEMF calculation.

From Fig.2 it is seen that the order parameter is also an oscillating function of the magnetic field and its value in general is the larger, the larger is the value of effective attraction W_0. At point of topological transition the order parameter has a local oscillation maximum. After the point of topological transition charge ordering is damaged, though very slowly. This damage can be explained by the fact that the kinetic energy of charge carriers in the layer plane in the presence of the Landau levels is higher than in the absence of these levels. From the thermodynamic considerations, the latter must resist charge ordering. However, even at $B = 0$ charge ordering is not complete, that is $\delta < 1$. It occurs due to the fact that at $\Delta \neq 0$ the interlayer motion of charge carriers attempts to smooth the nonuniformity of their distribution in the layers.

Calculation of the field dependence of thermoEMF and discussion of the results

In the calculation of thermoEMF we will proceed from the kinetic Boltzmann equation. Its application yields the following general formula for the longitudinal component of thermoEMF:

$$\alpha_{zz} \equiv \alpha = \frac{\sum\limits_{\beta} \tau_\beta v_{z\beta}^2 \dfrac{\partial f^0(\varepsilon_\beta)}{\partial T}}{e\sum\limits_{\beta} \tau_\beta v_{z\beta}^2 \dfrac{\partial f^0(\varepsilon_\beta)}{\partial \zeta}} \tag{6}$$

In this formula, e is electron charge modulus, $\beta \equiv (n, x)$, τ_β is relaxation time, $v_{z\beta}$ is longitudinal velocity of charge carriers, $f^0(\varepsilon_\beta)$ is the Fermi-Dirac distribution function. Summation over the numbers of the Landau levels in this formula can be followed out in the case when relaxation time depends only on x. It is noteworthy that more precise results in this situation are provided by the use of the Kubo formalism [18], however, in case of the fulfillment of the condition $\omega_c\tau >> 1$ an approach based on the Kubo formalism is equivalent to that based on the Boltzmann equation. Such equivalency

is due to the fact that at $\omega_c\tau >> 1$ an expansion of the Landau levels can be directly related to relaxation time, and a shift of energy levels due to scattering can be ignored. Moreover, we will consider that $l/a >> 1$ and $h/\tau\Delta >> 1$, which corresponds to pure and perfect layered crystal.

For simplicity, we will consider relaxation time to be stable. Taking into account that at low (helium) temperatures the dominant mechanism of charge carrier scattering is scattering on potential of randomly charged impurities, such approximation is correct for an open Fermi surface, though somewhat rough for a closed Fermi surface.

Calculation of thermoEMF by formula (6) with regard to dispersion law (2) yields the following expression for it:

$$\alpha = \frac{\pi\alpha_0 A}{B + C}. \tag{7}$$

In this formula, $\alpha_0 = k/e$, and dimensionless coefficients A, B, C are defined as follows:

$$A = \sum_{l=1}^{\infty}(-1)^l f_l^{th} \int_{-(\gamma-b)}^{\sqrt{w^2\delta^2+1}} \left|y^{-1}\right| \sqrt{\left(1+w^2\delta^2-y^2\right)\left(y^2-w^2\delta^2\right)} \sin\left[\pi l b^{-1}\left(\gamma-y\right)\right] dy, \tag{8}$$

$$B = 0.5 \int_{-(\gamma-b)}^{\sqrt{w^2\delta^2+1}} \left|y^{-1}\right| \sqrt{\left(1+w^2\delta^2-y^2\right)\left(y^2-w^2\delta^2\right)} dy, \tag{9}$$

$$C = \sum_{l=1}^{\infty}(-1)^l f_l^{\sigma} \int_{-(\gamma-b)}^{\sqrt{w^2\delta^2+1}} \left|y^{-1}\right| \sqrt{\left(1+w^2\delta^2-y^2\right)\left(y^2-w^2\delta^2\right)} \cos\left[\pi l b^{-1}\left(\gamma-y\right)\right] dy. \tag{10}$$

In these formulae, $\gamma = \zeta/\Delta$, $b = \mu^* B/\Delta$, $w = W_0/\Delta$. Besides:

$$f_l^{th} = \left[\mathrm{sh}\left(\pi^2 lkT/\mu^* B\right)\right]^{-1}\left[1-\left(\pi^2 lkT/\mu^* B\right)\mathrm{cth}\left(\pi^2 lkT/\mu^* B\right)\right], \tag{11}$$

$$f_l^{\sigma} = \frac{\pi^2 lkT/\mu^* B}{sh\left(\pi^2 lkT/\mu^* B\right)}. \tag{12}$$

If $\gamma - b \geq -w\delta$, the lower integration limit in formulae (8) – (10) should be replaced by $w\delta$. We also note that the reason for b subtraction from γ are physical circumstances mentioned in section 2.

The results of calculating thermoEMF in weak quantizing magnetic fields are presented in Fig. 3.

First of all, from Fig. 3 it is seen that in the ordered state the oscillations of thermoelectric coefficient are bioperiodic, whereas in the disordered state they are monoperiodic. The reason for this difference lies in the fact that in the disordered state the Fermi surface of a charge-ordered layered crystal within the open Brillouin zone has five extreme sections by planes perpendicular to magnetic field direction: three identical maximum by planes $k_z = 0$ and $k_z = \pm\pi/a$ and two minimum by planes $k_z = \pm\pi/2a$. The frequencies corresponding to these sections are as follows:

$$F_l = \frac{\pi l\left(\zeta + \sqrt{W_0^2\delta^2 + \Delta^2}\right)}{\mu^*}, \tag{13}$$

$$F_l' = \frac{\pi l\left(\zeta + W_0\delta\right)}{\mu^*}. \tag{14}$$

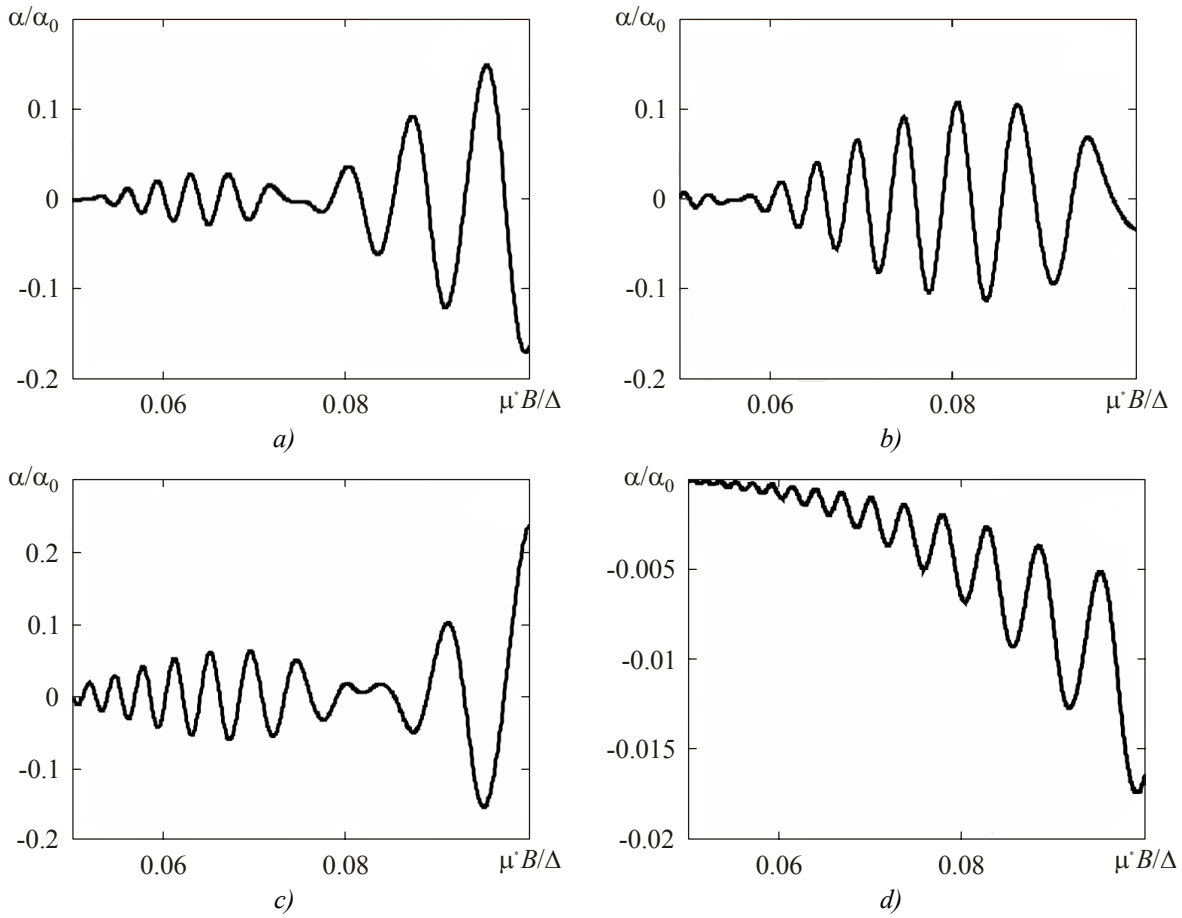

Fig. 3. Field dependence of thermoEMF in the range of magnetic fields $0.04 \le \mu^* B/\Delta \le 0.1$ at $\zeta_0/\Delta = 1$, $kT/\Delta = 0.03$ and: $W_0/\zeta_0 = 1.5$ (a), 2 (b), 2.5 (c), in the disordered state (d).

In a charge-ordered layered crystal these frequencies are close to each other, so high "carrier" frequencies are defined by the formula:

$$F_{lc} = \frac{\pi l\left(\zeta + W_0\delta\right)}{\mu^*}, \tag{15}$$

and low beat frequencies – by the formula:

$$F_{lb} = \frac{\pi l\Delta^*}{\mu^*}, \tag{16}$$

where $\Delta^* = 0.5\left(\sqrt{W_0^2\delta^2 + \Delta^2} - W_0\delta\right)$. At $\Delta = 0.01$ eV and $m^* = 0.01m_0$ and other parameters indicated in Fig. 3, the main high "carrier" frequencies in Figs. 3 *a*, *b*, *c*, respectively, are equal to 10.168, 10.154 and 10.547 T. Low beat frequencies are equal to 0.905, 0.678 and 0.542 T, respectively. In the disordered state, at selected problem parameters, the Fermi surface has the only extreme, namely maximum section by plane $k_z = 0$, so the basic high frequency equals 10.837 T and is modulated only by smooth component that cannot be taken into account within traditional approaches for which in a quasi-classical approximation the specific shape of the Fermi surface and its extent in the direction of a magnetic field are inessential.

From Fig. 3 it is also seen that charge ordering results in thermoEMF increase. Maximum values of thermoelectric coefficient in Figs. 3 *a*, *b*, *c* are equal to 12.9, 8.63 and 19.8 μV/K,

respectively. At the same time, in the disordered state (Fig. 3 *d*) maximum value of thermoEMF does not exceed 1.47 μV/K. Therefore, in the quasi-classical magnetic field region, at selected problem parameters, transition to charge-ordered state increases thermoEMF by a factor of 5.9 – 13.5. Moreover, in the charge-ordered state in the quasi-classical magnetic field region, there is thermoEMF polarity switching, whereas in the disordered state oscillations take place without polarity switching and thermoEMF remains purely negative, that is "electron" one. Within the assumed model of relaxation time this switching is attributable only to the transition from a closed Fermi surface into an open one with charge ordering.

The results of calculating thermoEMF in a wide magnetic field range are given in Fig. 4.

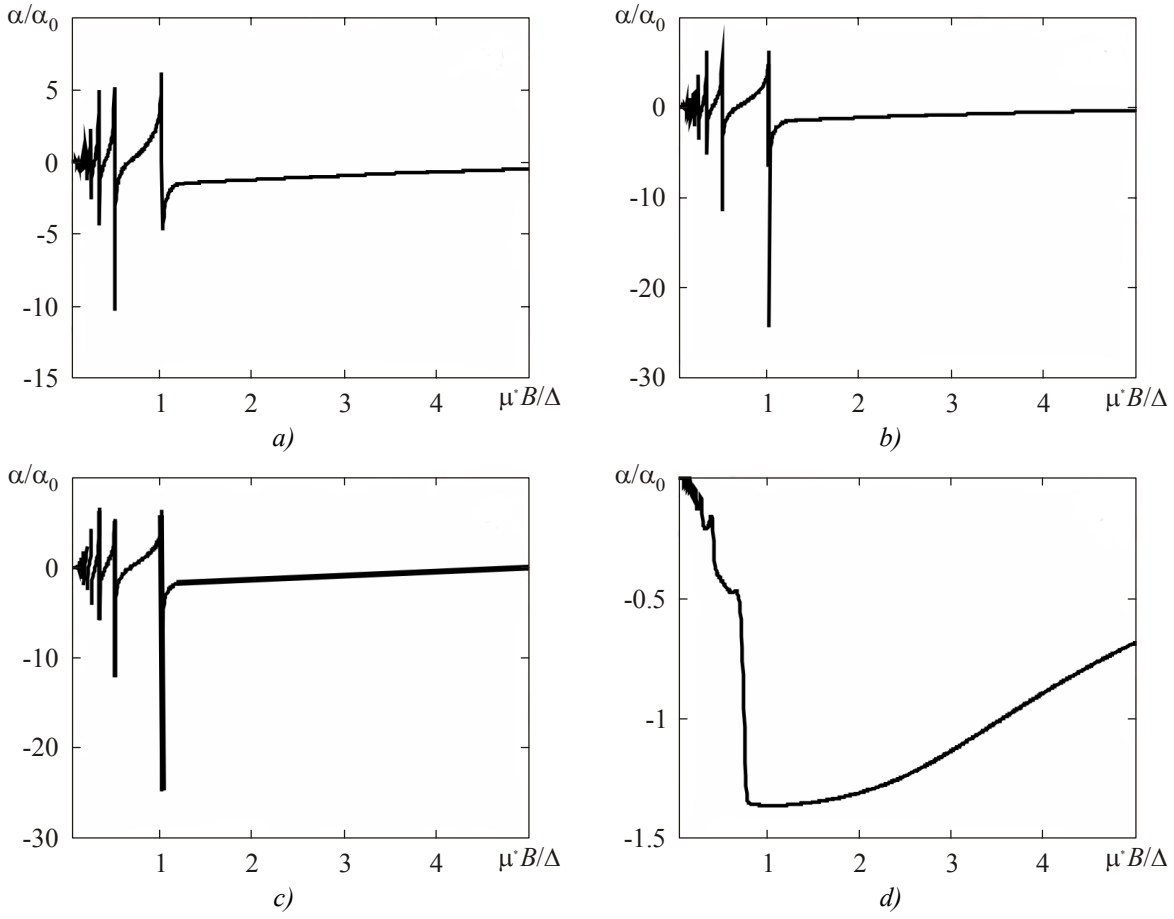

Fig. 4. Field dependence of thermoEMF in the magnetic field range $0.04 \le \mu^* B/\Delta \le 5$ *at* $\zeta_0/\Delta = 1$, $kT/\Delta = 0.03$ *and:* $W_0/\zeta_0 = 1.5$ *(a), 2 (b), 2.5 (c), in the disordered state (d).*

From Fig. 4 it is obvious that in the ordered state smooth oscillations of thermoEMF are gradually replaced by consecutive polarity switching (Figs. 4 *a*, *b*, *c*). This switching is synchronized with chemical potential changes up to the region of charge carrier condensation at the bottom of the Landau sub-band with the number $n = 0$. In so doing, each abrupt switching from positive to negative polarity corresponds to an abrupt hop of chemical potential, and a more smooth change from negative to positive polarity – to the region of a smooth change in the chemical potential. The last polarity switching occurs at point of topological transition, and thermoEMF reaches its maximum. ThermoEMF maxima in Figs. 4 *a*, *b*, *c* are equal to 890 μV/K, 2093 μV/K and 2147 μV/K, respectively. At the same time, in the disordered state (Fig. 4 *d*) thermoEMF almost everywhere remains negative and at point of the last oscillation maximum of the chemical potential it reaches maximum value equal to 117 μV/K, that is, by

a factor of 7.6 – 18.4 less than in the charge-ordered state. The existence of maximum thermoEMF value in a quantizing magnetic field can be explained by the following physical reasoning. On the one hand, it is clear that at low temperatures in a weak magnetic field the thermoEMF must tend to zero in conformity with general thermodynamic relations. On the other hand, it must also tend to zero in the ultra-quantum limit due to the Fermi surface compression. But it cannot be identical to zero. Hence, there must be the value of quantizing magnetic field induction whereby thermoEMF value reaches maximum. In our conditions it is induction from 0.86 to 1.73 T. However, with further increase in a magnetic field, the thermoEMF value starts decreasing to zero in the ordered state and more smoothly in the disordered state. A more drastic drop in thermoEMF value in the ordered state is explained by narrowing of conduction mini-band that describes the interlayer motion of charge carriers.

Finally, we establish an asymptotic law according to which thermoEMF is reduced to zero in the ultra-quantum limit. For this purpose at $\zeta_0 / \mu^* B << 1$ we write down relation (5) as follows:

$$\gamma = b - \sqrt{w^2\delta^2 + 1} + \frac{\pi^2 \zeta_0^2}{32\left(\mu^* B\right)^2 \sqrt{w^2\delta^2 + 1}}. \tag{17}$$

When calculating A, B, C coefficients in the integrals over y, let us replace the variables $y = y_1 - \sqrt{w^2\delta^2 + 1}$. In so doing, we take into account that on substituting (17) into expressions (8) – (10), in the trigonometric factors $(-1)^l$ will be compensated, following which cosines can be replaced by unities, and sines – by their arguments. The value f_l^{th} at $kT / \mu^* B << 1$ will be represented as:

$$f_l^{th} = \frac{\left(\pi^2 lkT / \mu^* B\right)^2}{2sh\left(\pi^2 lkT / \mu^* B\right)}, \tag{18}$$

Besides, we will take into account that at low η, as is shown by the numerical analysis, the following relations are valid:

$$\sum_{l=1}^{\infty} \frac{\eta l}{\mathrm{sh}\left(\eta l\right)} = \frac{2.467}{\eta}, \tag{20}$$

$$\sum_{l=1}^{\infty} \frac{\eta l^3}{\mathrm{sh}\left(\eta l\right)} = \frac{12.176}{\eta^3}. \tag{21}$$

We will expand the expressions under the integral signs in formulae (8) – (10) into a series in small parameter y_1 and make integration over it in the limits from 0 to $\dfrac{\pi^2 \zeta_0^2}{32\left(\mu^* B\right)^2 \sqrt{w^2\delta^2 + 1}}$.

Performing all operations and unification of numerical factors into one yields the following final expression for thermoEMF of charge-ordered layered crystal in the ultra-quantum limit:

$$\left|\alpha\right| = 0.305\alpha_0 \frac{\zeta_0^2 \Delta^2}{kT\left(\mu^* B\right)^2 \sqrt{W_0^2\delta^2 + \Delta^2}}. \tag{22}$$

At first sight, such a law is hard to be understood, since under this law thermoEMF does not tend to zero at $T = 0$. However, at real low temperatures this law does not lead to incorrect physical consequences. Indeed, for instance, at $B = 60$ T and with the above stipulated problem parameters,

taking into account the magnetic field dependence of the order, for values W_0 / ζ_0 equal to 1.5,2 and 2.5, respectively, we obtain thermoEMF values equal to 0.475, 0.358 and 0.287 μV/K, respectively, which is three orders less than maximum thermoEMF values. In the disordered state at $B = 60$ T thermoEMF is 0.725 μV/K, which is 162 times less than maximum value.

It is noteworthy that if for calculating thermoEMF in the ultra-quantum limit we use the same approach as in Ref. [19] for calculating conductivity, it will appear to be identical to zero, since the temperature dependence of distribution function in this approach is ignored, which actually corresponds to the case of $T = 0$. Thus, in any case this approach requires calculating thermoEMF in the following approximations by the small parameters kT / Δ, kT / ζ_0 and $kT / \mu^* B$.

Conclusions

1. It is shown that charge ordering leads to the appearance of a bioperiodic oscillation structure, increase of oscillation amplitude and switching of thermoEMF polarity in weak quantizing magnetic fields which is caused by the transition from a closed or transient Fermi surface into an open one and narrowing of conduction mini-band, describing the interlayer motion of charge carriers, when passing into a charge-ordered state.
2. In strong quantizing magnetic fields, charge ordering leads to an abrupt multiple switching of thermoEMF polarity and a sharp rise in its maximum value as compared to thermoEMF in the disordered state.
3. On reaching maximum value, the thermoEMF in both cases is drastically reduced by the law $\alpha \infty T^{-1} B^{-2}$, which is due to crystal Fermi surface compression in the direction of a magnetic field in the ultra-quantum limit. In a charge-ordered state, the thermoEMF is reduced even more drastically which is due to narrowing of a mini-band that describes the interlayer motion of charge carriers with charge ordering and a slow destruction of charge ordering in a magnetic field

It is clear that all the results obtained in the paper call for experimental validation, but all the experiments known to its author so far are concerned exceptionally with the research on the Shubnikov-de-Haas effect in layered crystals with strongly open Fermi surfaces in the field of application of a quasi-classical approximation.

The author is also aware that after reading this paper the readers may ask a question: what is the practical application, particularly in thermoelectricity, of the results obtained? None, if we speak about thermoelectric generators or coolers. And it can be very essential, if we speak about the thermoelectric effects as an instrument for studying other characteristics of material. Traditionally, until quite recently, the band structure of materials and the mechanisms of charge carrier scattering in them have been studied exceptionally with the aid of de Haas-van Alphen and Shubnikov-de Haas effects. However, it is difficult to single out de Haas-van Alphen and Shubnikov-de Haas oscillations from the general magnetic susceptibility and, the more so, conductivity. At the same time, thermoEMF oscillations are observed directly. So, their research at low temperatures can be considered as a further method for studying the band structure of materials, including thermoelectric ones, and the mechanisms of charge carrier scattering therein.

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