Quantum oscillations of the thermomagnetic coefficients of layered conductors in a strong magnetic field

O. V. Kirichenko and I. V. Kozlov

B. I. Verkin Institute for Low Temperature Physics and Engineering of the National Academy of Sciences of Ukraine, pr. Lenina 47, Kharkov 61103, Ukraine

D. Krstovska

Faculty of Natural Sciences and Mathematics, Physical Institute, P.O. Box 162, 1000 Skopje, Republic of Macedonia

V. G. Peschanskiĭ[a)]

B. I. Verkin Institute for Low Temperature Physics and Engineering of the National Academy of Sciences of Ukraine, pr. Lenina 47, Kharkov 61103, Ukraine

(Submitted February 8, 2008)
Fiz. Nizk. Temp. **34**, 681–686 (July 2008)

The linear response of the electronic system of a conductor to a perturbation in the form of an electric field and a temperature gradient in a quantizing magnetic field **B** is investigated theoretically. The thermoelectric effect in a layered conductor is analyzed and it is shown that the quasi-two-dimensional character of the dispersion law of the charge carriers results in gigantic oscillations of the thermo-emf. © *2008 American Institute of Physics*. [DOI: 10.1063/1.2957285]

I. INTRODUCTION

Landau's prediction that the magnetization metals oscillates as a function of the magnetic field strength[1] played an important role in solving the inverse problem of reconstructing the electronic energy spectrum of metals from the experimental data.[2,3] Kosevich's classical works devoted to the investigation of the quantum oscillations of the magnetic susceptibility and magnetoresistance of metals, performed together with Lifshits[3] and Andreev,[4,5] with the most general assumptions about the dispersion law for charge carriers made it possible to study in detail the form of the Fermi surface (FS) of practically all metals and subsequently also layered conductors. These works employed the area quantization rule to determine the quantized energy levels of conduction electrons undergoing finite motion in a plane orthogonal to the magnetic field:

$$S(\varepsilon,p_B)=\frac{2\pi\hbar eB}{c}\left(n+\frac{1}{2}\right),$$

where n are non-negative integers and $S(\varepsilon,p_B)$ is the area of the section of an isoenergy surface $\varepsilon(\mathbf{p})=\varepsilon$ by the plane $p_B=(\mathbf{p}\cdot\mathbf{B})/B=\text{const}$. The periods of oscillation of the magnetic susceptibility and the kinetic coefficients associated with this quantization are determined by the extremal values S_e of the area of the section of the FS. The contribution of each extremal section of the FS results in the appearance of harmonics of the form $\cos[kcS_e/2eB\hbar+(\pi/4)s]$, where $s=\text{sign}(\partial^2 S_e/\partial p_B^2)$.

The quantum oscillatory effects are most strongly manifested in layered conductors possessing substantial sections with sharply anisotropic, metal-type, electrical conductivity. The electrical conductivity in the plane of the layers is several orders of magnitude higher than that in a direction **n** normal to the layers; this suggests that the overlapping of the wave functions of the electrons belonging to different layers is small. In calculations of the kinetic coefficients, such an anisotropy can be described using a quasi-two-dimensional electronic energy spectrum, taking account of the fact that the energy $\varepsilon(\mathbf{p})$ of the conduction electrons depends weakly on the projection of their momentum $p_z=\mathbf{n}\cdot\mathbf{p}$ on the normal **n** to the layers. As a result, a substantially larger fraction of the charge carriers is drawn into the formation of the quantum oscillation effect than in ordinary metals. Investigations of the galvanomagnetic phenomena in many layered conductors at low temperatures have made it possible to study the topological structure of the Fermi surface $\varepsilon(\mathbf{p})=\varepsilon_F$, which is open. Specifically, there is every reason to believe that the FS of the tetrathiafulvalene salts $(\text{BEDT-TTF})_2\text{IBr}_2$ and $(\text{BEDT-TTF})_2\text{I}_3$ is a cylinder with weak fluting along the p_z axis.[6,7] If the magnetic field is strongly tilted away from the plane of the layers, then all flat sections of such a FS are closed, and all charge carriers in such conductors contribute to the quantum oscillations of the thermodynamic and kinetic properties in a strong magnetic field **B**.

The difference between the maximum $S_{\max}$ and minimum $S_{\min}$ sections of the FS in layered conductors with a quasi-two-dimensional electronic energy spectrum is small because of the weak fluting of the surface. Summing the contributions from different extremal sections, the oscillations acquire the form of beats $\cos[kc(S_{\max}+S_{\min})/2eB\hbar]\cos[kc(S_{\max}-S_{\min})/2eB\hbar-\pi/4]$.[8,9] Low-frequency oscillations with frequency proportional to $S_{\max}-S_{\min}$ are also observed in layered conductors; the amplitude of these oscillations is small but decays with temperature much more weakly than the amplitude of the Shubnikov-de Haas oscillations in ordinary metals. Slow oscillations of this type were first observed experimentally in measurements of the magnetoresistance of the compound β-$(\text{BEDT-TTF})_2\text{IBr}_2$.[10] They have now been observed in many organic metals.

In what follows, we shall examine the quantum oscillations of the kinetic coefficients in the presence of a temperature gradient. Specifically, we shall study the $1/B$ dependence in the form of gigantic oscillations of the thermoelectric field with amplitude much greater than the smooth part of the field. Analysis of the experimental results obtained for the oscillatory dependence of the thermo-emf on $1/B$ makes it possible to obtain detailed information about the energy spectrum of charge carriers in the conductor and is a very refined tool for studying the structure of this spectrum.

II. LINEAR RESPONSE OF THE ELECTRONIC SYSTEM TO A PERTURBATION IN THE FORM OF AN ELECTRIC FIELD AND A TEMPERATURE GRADIENT

The electric current density $\mathbf{j}$ and heat flux density $\mathbf{q}$ which arise in a conductor as a result of an external perturbation in the form of a temperature gradient ∇T and an electric field $\mathbf{E}$ have the form

$$j_i = \sigma_{ij} E_j^* - \alpha_{ij} \frac{\partial T}{\partial x_j}, \tag{1}$$

$$q_i = \beta_{ij} E_j^* - \kappa_{ij} \frac{\partial T}{\partial x_j}, \tag{2}$$

where

$$E_j^* = E_j - \frac{1}{e} \frac{\partial \mu}{\partial x_j}, \tag{3}$$

and μ is the chemical potential of the electrons.

The construction of a linear theory of thermomagnetic phenomena reduces to calculating the kinetic coefficients $\sigma_{ij}(\mathbf{B})$, $\alpha_{ij}(\mathbf{B})$, $\beta_{ij}(\mathbf{B})$, and $\kappa_{ij}(\mathbf{B})$ which relate the fluxes with small perturbations of the electronic system. We shall assume that the electron-electron interaction establishes in a conductor, over a time much shorter than the damping time of macroscopic fluxes, a local quasi-equilibrium distribution of charge carriers with coordinate-dependent parameters (T, μ). We shall examine the case of low temperatures, where the number of phonons is small and the main relaxation mechanism is elastic scattering of electrons by impurity centers, whose concentration is not too large so that the revolution frequency ω_c of the charge carriers in a magnetic field is much higher than the carrier-scatterer collision frequency.

In the absence of current-conducting contacts, a temperature gradient generates a thermoelectric field in the conductor. Setting $\mathbf{j}=0$ in Eq. (1) we obtain

$$E_i = \rho_{il} \alpha_{lj} \frac{\partial T}{\partial x_j} + \frac{1}{e} \frac{\partial \mu}{\partial x_i}, \tag{4}$$

where ρ_{ij} is the resistivity tensor, which is the inverse of the conductivity tensor σ_{ij}.

The gradient of the chemical potential is easily found from the condition that the number of charge carriers per unit volume is constant:

$$\frac{\partial N}{\partial \mathbf{r}} = 0, \tag{5}$$

where

$$N = \frac{2eB}{c(2\pi\hbar)^2} \sum_{n=0}^{\infty} \int \frac{dp_H}{1 + \exp\{[\varepsilon_n(p_B) - \mu_\pm]/T\}}. \tag{6}$$

Applying Poisson's relation

$$\sum_{n=0}^{\infty} \phi_n = \int_{-1/2}^{\infty} dn \phi(n) \sum_{k=-\infty}^{\infty} \exp(2\pi i k n)$$

to Eq. (5) and switching from integration over n to integration over energy, we obtain in the leading approximation in the small parameter $\hbar\omega_c/\mu$

$$\nabla\mu = -\nabla T \frac{\pi^2 T}{3\nu(\mu)} \left[\frac{\partial \nu(\mu)}{\partial \mu} - \frac{2^{3/2}}{\hbar^3 \pi^{1/2}} \sum_{k=1}^{\infty} (-1)^k k^{1/2} \times \sum_e P_k(u) \frac{m^{3/2}}{(\hbar\omega_c)^{1/2}} \left| \frac{\partial^2 S_e}{\partial p_H^2} \right|^{-1/2} \sin\left(\frac{kcS_e}{e\hbar H} + \frac{\pi}{4} s \right) \right], \tag{7}$$

where $\upsilon(\varepsilon)$ is the electron density of states in the absence of a quantizing magnetic field, $m = (2\pi)^{-1} \partial S/\partial\varepsilon$ is the effective mass of the conduction electrons, $u = 2\pi^2 T/\hbar\omega_c$, and the function

$$P_k(u) = -\frac{3}{ku} \frac{\sinh(ku) - ku\cosh(ku)}{\sinh^2(ku)}$$

goes to 1 as $T \to 0$. Since charge carriers on the extremal sections of the FS form the oscillations, the summation in Eq. (7) must extend over the electron states in all of these sections.

In a quasi-isotropic conductor at sufficiently low temperature, the amplitude of the oscillating terms in Eq. (7) is greater than the first term by at least a factor of $(\mu/\hbar\omega_c)^{1/2}$.

The amplitudes of the quantum oscillations of both terms on the right-hand side of Eq. (3) are of the same order of magnitude, and to calculate the thermoelectric field it is very important to take account of the oscillations of $\nabla\mu$. Ordinarily, when calculating the macroscopic fluxes in the presence of a temperature gradient, the contributions associated with the magnetization $\mathbf{M}$ of the electronic system are excluded from them. It is assumed that the conduction current density differs from the average microscopic current density $\mathrm{Tr}(e\hat{f}\hat{\mathbf{v}})$ by the vector c curl $\mathbf{M}$, and the energy flux minus the magnetic energy flux is studied as the heat flux ($\hat{\mathbf{v}}$ and e are the velocity operator and the electron charge, $\hat{f}$ is the density operator, and c is the speed of light). Eliminating the "nonthermal" contribution in this manner from the heat flux removes the discrepancy between the computational results and Onsager's principle, which follows from the condition that the entropy attains its maximum value in the equilibrium state. The effects associated with the magnetism of the conduction electrons must be taken into account when calculating the off-diagonal components of the kinetic coefficients in a quantizing magnetic field.[11]

For elastic scattering, the diagonal components of the thermomagnetic coefficients in the single-electron approximation are related with the diagonal components of the conductivity tensor. The components σ_{ii} can be calculated using Kubo's relation[12]

$$\sigma_{ii}(T,\mu)=\int\left(-\frac{\partial f_0(E)}{\partial E}\right)F_{ii}(E)dE, \tag{8}$$

where

$$F_{ii}(E)=\pi\hbar e^2\mathrm{Sp}\langle\delta(E-\hat{H})\hat{v}_i\delta(E-\hat{H})\hat{v}_i\rangle. \tag{9}$$

Here $f_0(E)$ is the Fermi distribution function and the angular brackets denote averaging over the configurations of randomly arranged impurity centers. The components α_{ii}, β_{ii}, and κ_{ii} satisfy[13]

$$\beta_{ii}(T,\mu)=T\alpha_{ii}(T,\mu)=\int\left(-\frac{\partial f_0(E)}{\partial E}\right)\frac{E-\mu}{e}F_{ii}(E)dE, \tag{10}$$

$$\kappa_{ii}(T,\mu)=\int\left(-\frac{\partial f_0(E)}{\partial E}\right)\frac{(E-\mu)^2}{e^2T}F_{ii}(E)dE. \tag{11}$$

Thus, all diagonal components of the electronic kinetic coefficients can be expressed in terms of the function $F_{ii}(E)$.

III. THERMOELECTRIC EFFECT IN A LAYERED CONDUCTOR

We shall now examine the thermoelectric effect in a layered conductor when the temperature gradient and the magnetic field $\mathbf{B}=(0,0,B)$ are directed along the normal to the layers. It is easy to show that in this case the thermoelectric field is also directed transverse to the layers

$$E_z=\frac{\alpha_{zz}}{\sigma_{zz}}\frac{\partial T}{\partial z}+\frac{1}{e}\frac{\partial\mu}{\partial z} \tag{12}$$

and is determined only by the diagonal matrix elements of the kinetic coefficients in the asymptotic approximation with respect to the small parameter characterizing the quasi-two-dimensionality of the energy spectrum of the conduction electrons. Many works use a quite simple model of the quasi-two-dimensional dispersion law for charge carriers to interpret the experimentally studied phenomena in organic conductors:[14,15]

$$\varepsilon(\mathbf{p})=\frac{p_x^2+p_y^2}{2m}-2t\cos\frac{ap_z}{\hbar}, \tag{13}$$

where $m=$const, a is the distance between the layers, and t is determined by the overlap integral of the wave functions of electrons from neighboring layers and is much less than the Fermi energy but greater than the splitting $\hbar\omega_c$ between the quantized energy levels in currently attainable magnetic fields. As will be shown below, at temperatures much higher than the splitting between the quantized energy levels of the electrons, in the case of the dispersion law (13) the thermoelectric effect transverse to the layers, when the temperature gradient and the magnetic field are parallel to the normal to the layers, is negligibly small, and even vanishes in the absence of a magnetic field.

For elastic scattering by randomly distributed impurity centers we shall write the Hamiltonian of an electron in the form

$$\hat{H}=\hat{\varepsilon}+\sum_i\hat{V}_i, \tag{14}$$

where $\hat{V}_i=\hat{V}(\hat{\mathbf{r}}-\mathbf{R}_i)$ is the potential of an impurity center located at the point $\mathbf{R}_i$. We assume that the scattering potential is weak and that its range is the smallest parameter with the dimension of length in the problem.

The function F_{ii} must be determined in order to use Eqs. (3)–(5) to determine the thermomagnetic coefficients. The operator $\delta(E-\hat{\boldsymbol{H}})$ in the expression (9) for F_{ii} can be represented in the form

$$\delta(E-\hat{H})=\frac{i}{2\pi}[\hat{G}^+(E)-\hat{G}^-(E)], \tag{15}$$

where $\hat{\boldsymbol{G}}^{\pm}(E)=(E-\hat{\boldsymbol{H}}\pm i\delta)^{-1}$—the single-electron Green's function.

In the self-consistent Born approximation, the Green's function can be "de-coupled" when averaging over the impurities[16]

$$\langle\hat{G}^{\pm}\hat{v}_i\hat{G}_{\pm}\hat{v}_j\rangle=\langle\hat{G}^{\pm}\rangle\hat{v}_i\langle\hat{G}_{\pm}\rangle\hat{v}_j. \tag{16}$$

Then the Green's $\hat{\boldsymbol{G}}^{\pm}(E)$ assumes the form

$$\langle\hat{G}^{\pm}(E)\rangle=\frac{1}{E-\hat{\varepsilon}-\hat{\Sigma}^{\pm}(E)}, \tag{17}$$

where $\hat{\boldsymbol{\Sigma}}^{\pm}(E)=\langle\hat{\boldsymbol{\Sigma}}_i^{\pm}(E)\rangle$ is the self-energy part averaged over all impurity centers (see Refs. 16 and 17),

$$\hat{\Sigma}_i^{\pm}(E)=\hat{V}_i+\hat{V}_i\langle\hat{G}^{\pm}(E)\rangle\hat{V}_i+\ \dots\ . \tag{18}$$

As a result of averaging, the operator $\hat{\Sigma}^{\pm}(E)$ becomes diagonal and can be represented in the form $\hat{\Sigma}^{\pm}(E)=\hat{\Sigma}^{\pm}(E)\hat{I}$, where $\hat{I}$ Is the unit operator.

It is easy to see that the quantity $\hat{\Sigma}^{\pm}(E)$ is related with the scattering tensor

$$\hat{T}_i^{\pm}(E)=\hat{V}_i+\hat{V}_i\hat{G}_0^{\pm}(E)\hat{V}_i+\dots \tag{19}$$

by the relation

$$\hat{\Sigma}^{\pm}(E)=\langle\hat{T}^{\pm}(E-\Sigma^{\pm}(E))\rangle. \tag{20}$$

The scattering tensor for a conductor with the dispersion law (13) is calculated in Ref. 9, where the method developed in Refs. 5 and 18 is used. Following these works, we shall represent the Green's function neglecting impurity scattering $\hat{G}_0(E)=(E-\hat{\varepsilon}\pm i\delta)^{-1}$ in the coordinate representation in the form of two terms:

$$G_0^{\pm}(\mathbf{r},\mathbf{r}';E)=\Phi(\mathbf{r},\mathbf{r}')[G_{\mathrm{cl}}(\mathbf{r}-\mathbf{r}';E)+G_q^{\pm}(E)], \tag{21}$$

where the factor $\Phi(\mathbf{r},\mathbf{r}')$ depends on the gauge of the vector potential $\mathbf{A}$. In the Landau gauge $\mathbf{A}=(0,Bx,0)$

$$\Phi(\mathbf{r},\mathbf{r}')=\exp\left[\frac{i\hbar c}{2eB}(x+x')(y-y')\right].$$

Here G_{cl} is the real part of the Green's function $G_0^{\pm}$ in the absence of a magnetic field, and the coordinate dependence of $G_q^{\pm}$ can be neglected for a short-range impurity, which is the case being considered here. The function $G_{\mathrm{cl}}(\mathbf{r},\mathbf{r}';E)$ appears only in the expression for the total scattering ampli-

tude, which is related with the impurity potential by the relation

$$f_{\text{imp}} = \frac{m}{2\pi\hbar^2}\int V(\mathbf{r})\psi_0(\mathbf{r})d^3r,$$

$$\psi_0(\mathbf{r}) = 1 + \int G_{\text{cl}}(\mathbf{r},\mathbf{r}';E)V(\mathbf{r}')\psi_0(\mathbf{r}')d^3r', \tag{22}$$

where, since G_{cl} is a weak function of E, we can set $E \approx \mu$. The result for the scattering tensor is the expression

$$\langle T^{\pm}(E)\rangle = \frac{\dfrac{2\pi\hbar^2}{m}f_{\text{imp}}n_{\text{imp}}}{1 - \dfrac{2\pi\hbar^2}{m}f_{\text{imp}}G_q^{\pm}(E)}, \tag{23}$$

where n_{imp} is the impurity concentration,

$$G_q^{\pm}(E) = \mp\frac{im}{2\hbar^2 a}\left[1 + 2\sum_{k=1}^{\infty}(-1)^k \times\exp\left(\pm\frac{2\pi ikE}{\hbar\Omega}\right)J_0\left(\frac{4\pi kt}{\hbar\Omega}\right)\right], \tag{24}$$

and J_n is a Bessel function. We note that when Eq. (20) is solved in the leading approximation in the small parameter $\hbar\omega_c/t \ll 1$ the oscillating part $\Sigma^{\pm}(E)$ in the argument of the scattering tensor can be neglected, and since the total amplitude of scattering by an impurity is small, only the contribution linear in $G_q^{\pm}$ need be retained in the expression for the scattering tensor (23).

The expressions (23) and (24), together with the relation (20), make it possible to find the Green's function $\langle\hat{G}^{\pm}\rangle$ from the relation (17). Substituting into Eq. (9), it is easy to see that $\Sigma^{\pm}(E)$ enters into the expression (9) in the form of the combination $(i/h)[\Sigma^{+}(E)-\Sigma^{-}(E)] = 1/\tau(E)$, which is the reciprocal of the relaxation time. As a result, we obtain the following expression for the function F_{zz}:

$$F_{zz}(E) = A\tau(E)\times\left[1 + \sum_{k=1}^{\infty}\left(\frac{\hbar\omega_c}{\pi k} + \frac{\hbar}{\tau(E)}\right)\frac{(-1)^k}{t}D_k\sin\left(\frac{2\pi k\tilde{E}}{\hbar\omega_c}\right)J_1\left(\frac{4\pi kt}{\hbar\omega_c}\right)\right]. \tag{25}$$

Here

$$\frac{1}{\tau(E)} = \frac{1}{\tau_0}\left[1 + 2\sum_{k=1}^{\infty}(-1)^k D_k\cos\left(\frac{2\pi k\tilde{E}}{\hbar\omega_c}\right)J_0\left(\frac{4\pi kt}{\hbar\omega_c}\right)\right], \tag{26}$$

$$\tilde{E} = E - \text{Re}\,\Sigma_{\text{cl}}(\mu), \quad A = \frac{n_e e^2(2t)^2 a^2}{2\mu\hbar^2},$$

$$\frac{1}{\tau_0} = \frac{4\pi^2\hbar n_{\text{imp}}}{ma}f_{\text{imp}}^2, \quad D_k = \exp\left(-\frac{\pi k}{\omega_c\tau_0}\right),$$

and n_e is the density of the charge carriers. The conductivity component σ_{zz} calculated using the relation for F_{zz} and the relation (8) is identical to the expression obtained for σ_{zz} and analyzed in detail by Grigor'ev.[19] The asymptotic behavior of σ_{zz} for $(\hbar\omega_c/t) \ll 1$ is given by the expression

$$\sigma_{zz} = \sigma_0\left\{1 + 2D_1R_1\sqrt{\frac{\hbar\omega_c}{2\pi^2 t}}\cos\left(\frac{2\pi\mu}{\hbar\omega_c}\right)\cos\left(\frac{4\pi t}{\hbar\omega_c} - \frac{\pi}{4}\right) + D_1^2\frac{\hbar\omega_c}{2\pi^2 t}\cos\left[2\left(\frac{4\pi t}{\hbar\omega_c} - \frac{\pi}{4}\right)\right]\right\}. \tag{27}$$

Here we have retained only the first term of the series with respect to k, $\sigma_0 = A\tau_0$, $R_k = ku/\sinh(ku)$. At low temperatures ($u \ll 1$) the smooth part of σ_{zz} is at least $(t/\hbar\omega_c)^{1/2}$ times greater than the oscillating correction. The third term in braces, which describes slow oscillations, can dominate at higher temperatures, since it does not contain the temperature factor R_1.

Substituting F_{zz} into the expression (10) gives the following expression for α_{zz}:

$$\alpha_{zz} = \frac{\sigma_0}{e}\frac{4\pi^3}{3}\frac{T}{\hbar\omega_c}\sum_{k=1}^{\infty}(-1)^k kD_k\sin\left(\frac{2\pi k\mu}{\hbar\omega_c}\right)P_k(u) \times\left\{J_0\left(\frac{4\pi kt}{\hbar\omega_c}\right) - \left(\frac{1}{\pi k} + \frac{1}{\omega_c\tau_0}\right)\frac{\hbar\omega_c}{2t}J_1\left(\frac{4\pi kt}{\hbar\omega_c}\right)\right\}. \tag{28}$$

Using the asymptotic expansion for Bessel functions, we obtain for the first harmonic

$$\alpha_{zz} = -\frac{2\pi^2 T}{3e}\sigma_0 D_1 P_1(u)\sqrt{\frac{2}{\hbar\omega_c t}}\sin\left(\frac{2\pi\mu}{\hbar\omega_c}\right) \times\cos\left(\frac{4\pi t}{\hbar\omega_c} - \frac{\pi}{4}\right). \tag{29}$$

The coefficient α_{zz} does not have a component that varies smoothly as a function of B because the "classical" part of the density of states of the electrons in the case of the dispersion law (13) does not depend on the energy.

Therefore the thermoelectric field E^*_z can be written in the form

$$E_z^* = \frac{\alpha_{zz}}{\sigma_0}\frac{\partial T}{\partial z}. \tag{30}$$

In the expression (30) we have neglected the terms which arise as a result of the interference of quantum oscillations of σ_{zz} and α_{zz}, whose amplitude is small in the parameter $\hbar\Omega/\mu \ll 1$ as compared with the leading oscillatory contribution.

In the model considered here for the dispersion law for charge carriers, it is easy to find the gradient of the chemical potential as

$$\frac{\partial\mu}{\partial z} = \frac{4\pi^3}{3}\frac{T}{\hbar\omega_c} \times\sum_{k=1}^{\infty}(-1)^k kD_k P_k(u)\sin\left(\frac{2\pi k\mu}{\hbar\Omega}\right)J_0\left(\frac{4\pi kt}{\hbar\omega_c}\right)\frac{\partial T}{\partial z} \tag{31}$$

and to show that the leading approximation in the small parameter $\hbar\omega_c/t$ the quantity $\partial\mu/\partial z$ equals eE^*_z.

When the magnetic field is tilted with respect to the normal to the layers by the angle ϑ the drift velocity of the charge carriers along the magnetic field

$$\bar{v}_z = \frac{at}{\hbar} J_0\left(\frac{ap_\perp}{\hbar}\tan\vartheta\right)\sin\left(\frac{ap_B}{\hbar\tan\vartheta}\right) \tag{32}$$

depends on $p_\perp = (2m\varepsilon)^{1/2}$ and the part $E_z^{(\mathrm{mon})}$ of the thermoelectric field that varies monotonically as a function of B vanishes in the leading approximation in the small parameter $\eta = t/\mu$ characterizing the quasi-two-dimensionality of the energy spectrum only for distinguished values of the angle ϑ corresponding to the zeros of the vessel function $J_n(ap_\perp \tan\vartheta/\hbar)$ with $n=0$, 1. In all other cases the monotonic part of the coefficient α_{zz} is different from zero and the amplitude of the oscillations of the thermoelectric field is $(\mu/\eta\hbar\omega_c)^{1/2}$ times greater than its monotonic part.

For a model more complicated than the model (13) for the quasi-two-dimensional energy spectrum of charge carriers, the monotonic part of the thermoelectric field is different from zero even in a magnetic field directed along the normal to the layers. Then the amplitude of the oscillations is, once again, at least $(\mu/\eta\hbar\omega_c)^{1/2}$ times greater than $E_z^{(\mathrm{mon})}$, where the quasi-two-dimensionality parameter η determines the magnitude of the fluting of the FS. In organic layered conductors, μ is ordinarily of the order of 0.1 eV, while the splitting $\hbar\omega_c$ between the Landau levels in real magnetic fields does not exceed 1 meV.[6] Since the parameter η is small, $(\mu/\eta\hbar\omega_c)^{1/2}$ can be expected to be of the order of 10^2.

IV. CONCLUSIONS

In layered conductors with a quasi-two-dimensional dispersion law for charge carriers, the dependence of the thermoelectric field on the reciprocal of the magnetic field exhibits gigantic oscillations and contains rich information about the charge carriers. In contrast to galvanomagnetic phenomena, thermoelectric effects are much more sensitive to the choice of model for the electronic energy spectrum. The simple FS model (13) widely used for a quasi-two-dimensional conductor to analyze experimental results is "exotic," since in this model the monotonic part of E_z vanishes at $\vartheta=0$ and the thermoelectric effect is absent irrespective of the quantization conditions for the orbital motion of conduction electrons. However, if the FS is a weakly fluted cylinder of arbitrary form, then the dependence of the thermoelectric field E_z on the reciprocal of the magnetic field, once again, exhibits oscillations whose amplitude is much greater than the monotonic part of E_z.

Experimental investigations of the ratios of the phases of the oscillations of the thermoelectric coefficients make it possible to determine how closely the model (13) corresponds to the real dispersion law for charge carriers in a conductor. The large amplitude of the oscillations makes it possible to determine the extremal sections of the FS of layered conductors more accurately. The results of experimental studies of the thermoelectric effect together with data from galvanomagnetic measurements make it possible to determine the effective masses of the charge carriers drawn into the oscillations as well as the parameter characterizing the quasi-two-dimensionality of the electronic spectrum.

We thank the Ministry of Education and Science of the Republic of Macedonia for providing financial support for this work.

[a]E-mail: vpeschansky@ilt.kharkov.ua

[1]L. D. Landau, Z. Phys. **64**, 629 (1930).
[2]L. Onsager, Philos. Mag. **43**, 1006 (1952).
[3]I. M. Lifshits and A. M. Kosevich, Zh. Éksp. Teor. Fiz. **29**, 730 (1955) [Sov. Phys. JETP **29**, 636 (1955)].
[4]A. M. Kosevich and V. V. Andreev, Zh. Éksp. Teor. Fiz. **38**, 882 (1960) [Sov. Phys. JETP **38**, 637 (1960)].
[5]V. V. Andreev and A. M. Kosevich, Zh. Éksp. Teor. Fiz. **43**, 1061 (1962) [Sov. Phys. JETP **15**, 735 (1962)].
[6]M. V. Kartsovnik, Chem. Rev. (Washington, D.C.) **104**, 5737 (2004).
[7]M. V. Kartsovnik and V. G. Peschanskiĭ, Fiz,. Nizk. Temp. **31**, 249 (2005) [Low Temp. Phys. **31**, 185 (2005)].
[8]P. D. Grigoriev, M. V. Kartsovnik, W. Biberacher, N. D. Kusch, and P. Wider, Phys. Rev. B **65**, 060403 (2002).
[9]O. V. Kirichenko and I. V. Kozlov, Fiz,. Nizk. Temp. **28**, 509 (2002) [Low Temp. Phys. **28**, 359 (2002)].
[10]M. V. Kartsovnik, V. Laukhin, V. Nizhankovskiĭ, and A. Ignat'ev, Pis'ma Zh. Éksp. Teor. Fiz. **48**, 187 (1965).
[11]Yu. N. Obraztsov, Fiz. Tverd. Tela **6**, 414 (1964) [Sov. Phys. Solid State **6**, 331 (1964)].
[12]G. Kubo, J. Phys. Soc. Jpn. **12**, 570 (1957).
[13]V. G. Bar'yakhtar and S. V. Peletminskiĭ, Zh. Éksp. Teor. Fiz. **48**, 187 (1965)] [JETP **21**, 126 (1965)].
[14]J. Wosnitza, *Fermi Surfaces of Low-Dimensional Organic Metals and Superconductors*, Springer Tracks in Modern Physics (1996).
[15]J. Singleton, *Studies of Quasi-Two-Dimensional Organic Conductors Based on BEDT-TTF Using High Magnetic Fields*, Reports on Progress in Physics (2000).
[16]A. A. Abrikosov, L. P. Gor'kov, and I. E. Dzyaloshinskiĭ, *Methods of Quantum Field Theory in Statistical Physics*, Fizmatgiz, Moscow (1962).
[17]S. Doniach and E. H. Sondheimer, *Green's Functions for Solid-State Physicists*, Imperial College Press, London (1998).
[18]V. G. Skobov, Pis'ma Zh. Éksp. Teor. Fiz. **39**, 689 (1960).
[19]P. D. Grigoriev, Phys. Rev. B **67**, 144401 (2003).

Translated by M. E. Alferieff