ISSN 1063-7761, Journal of Experimental and Theoretical Physics, 2019, Vol. 129, No. 2, pp. 292–297. © Pleiades Publishing, Inc., 2019.
Russian Text © The Author(s), 2019, published in Zhurnal Eksperimental'noi i Teoreticheskoi Fiziki, 2019, Vol. 156, No. 2, pp. 348–354.

Shubnikov–de Haas Thermoelectric Field Oscillations in Layered Conductors in the Vicinity of a Topological Lifshits Transition

V. G. Peschansky[a,b,*], O. Galbova[c], and K. Yiasemides[d]

[a] *Karazin Kharkov National University, Kharkov, 61022 Ukraine*

[b] *Verkin Institute for Low Temperature Physics and Engineering, National Academy of Sciences of Ukraine, Kharkov, 61103 Ukraine*

[c] *Institute of Physics, Faculty of Natural Sciences and Mathematics, Ss. Cyril and Methodius University in Skopje, Skopje-1000, Republic of Macedonia*

[d] *University of West Attica, Egaleo, 12244 Greece*

**e-mail: peschansky@ilt.kharkov.ua*

Received January 23, 2019; revised March 15, 2019; accepted March 22, 2019

Abstract—We have studied the response of a layered conductor with a quasi-two-dimensional electron spectrum and a multisheet Fermi surface (FS) formed by two weakly corrugated cylinders and two planar sheets adjoining these cylinders to nonuniform heating. As a result of action of pressure on the conductor or its doping with impurity atoms, distance Δ_p between the FS cylinder and planar sheets can be reduced to such an extent that conduction electrons begin roaming over these sheets, tunneling from one FS sheet (cavity) to the other. If a conduction electron can visit several times all FS sheets during its mean free time, its motion in the plane orthogonal to the magnetic field becomes finite. In this case, Shubnikov–de Haas oscillations are excited with a period determined by the closed area circumscribed by the electron during its motion in the magnetic-breakdown trajectory in the momentum space. We have calculated the dependence of the thermoelectric field on the magnitude and orientation of a quite strong quantizing magnetic field. In a magnetic field normal to the layers, the cross sections of the cylindrical part of the FS are equidistant from both FS planar sheets. However, this equidistance is violated even for a small deviation of the field from the normal to the layers by angle ϑ, and, at a certain value ϑ_k of this angle, the probability of magnetic breakdown to one of the FS planar sheets can be so low that the electron cannot close the magnetic-breakdown trajectory, and its motion over the other planar sheet with a visit to the FS cylindrical part becomes infinite. In this case, the magnetic-breakdown quantum oscillations of magnetization and all kinetic characteristics of the conductor disappear, and their disappearance is repeated periodically with a change in the slope of the magnetic field to the layers as a function of $\tan\vartheta$.

DOI: 10.1134/S1063776119080077

1. INTRODUCTION

The electron topological transition in degenerate conductors predicted by Lifshits [1], during which the connectivity of the Fermi surface (FS) changes because of an external action in the form of pressure on such conductors or their doping with impurity atoms was soon detected in many metals and their alloys in the normal and superconducting states [2]. During the last two decades, the interest in the Lifshits topological transitions was shifted to the region of low-dimensional conductors. The topological structure of the FS in such conductors even under a low pressure can naturally change, which is accompanied by anomalies in the kinetic and thermodynamic electron characteristics.

Such anomalies are manifested most clearly in thermoelectric phenomena and contain important information on conduction electrons, even in classically strong magnetic fields **B**, when frequency ω_c of circumvention of a charge is much larger than frequency $1/\tau$ of electron collisions, and separation $\hbar\omega_c$ between quantized energy levels is much smaller than blurring temperature T of their equilibrium Fermi distribution function. Theoretical analysis of thermoelectric effects in layered conductors with a multisheet FS in this range of magnetic fields for different orientations of the temperature gradient and vector **B** [3, 4] paves the way to detailed investigation of the electron energy spectrum for such conductors, in particular, the determination of fine features of individual FS cavities (sheets) and their mutual arrangement in the momentum space; analysis of the temperature depen-

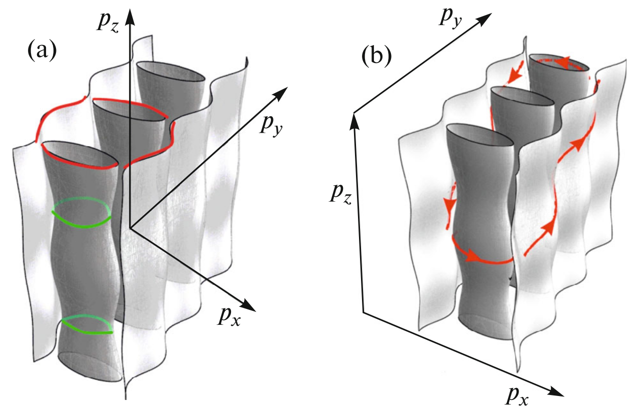

Fig. 1. (Color online) Electron trajectories on the FS in a layered conductor, which consists of cylinders in each cell of the momentum space and two adjoining quasi-planar sheets weakly corrugated along axis $p_z = \mathbf{p}\cdot\mathbf{n}$ (**n** is the normal to the layers): (a) magnetic-breakdown (red) and conventional (green) electron trajectories in a magnetic field normal to the layers; (b) magnetic-breakdown trajectories (red) in a field tilted to the layers.

dence of the thermoelectric coefficients makes it possible to study various mechanisms of charge carrier relaxation.

In stronger magnetic fields (for $\hbar\omega_c \geq 2\pi^2 T$), it is necessary to consider the energy quantization for charge carriers performing finite motion in the plane orthogonal to the magnetic field, which leads to an oscillatory dependence of the thermoelectric field on the reciprocal magnetic field. The oscillatory dependence of the magnetic susceptibility of metals in a quantizing magnetic field on the reciprocal value of the quantizing magnetic field was predicted by Landau [5] and was detected in the same year of 1930 at the Leiden laboratory in quite perfect bismuth single crystals at the liquid helium temperature [6]; later, this dependence was observed for the magnetoresistance of the same bismuth samples [7]. Theoretical calculation of magnetic susceptibility of metals in a quantizing magnetic field in the case of an isotropic dependence of energy $\varepsilon(\mathbf{p})$ of charge carriers on their momentum **p** was reported by Landau later in Appendix to the Schoenberg paper [8]. The magnetization of metals under the most general assumptions concerning the form of their electron energy spectrum was analyzed theoretically by Lifshits and Kosevich [9] and served as the beginning of a new trend of the electron physics of metals, which was later called fermiology, involving the solution of the inverse problem of reconstructing the FS shape from the measured oscillatory dependence of magnetization and kinetic coefficients on the reciprocal value of a strong magnetic field.

In this paper, we consider the response of the electron system of layered conductors with a quasi-two-dimensional electron energy spectrum of an arbitrary form to nonuniform heating near the Lifshits topological transition and analyze the dependence of thermoelectric field

$$E_i = \frac{\pi^2 T}{3e}\rho_{ik}\frac{\partial\sigma'_{kj}(\mu)}{\partial\mu}\frac{\partial T}{\partial x_j}, \qquad (1)$$

generated by temperature gradient $\partial T/\partial\mathbf{r}$ on the magnitude and orientation of the quantizing magnetic field. Here, ρ_{ik} is the resistivity tensor, which is reciprocal of conductivity tensor $\sigma_{ij}(\mu)$, and $\sigma'_{ij}(\mu)$ coincides with tensor $\sigma_{ij}(\mu)$ if we replace relaxation time τ of charge carriers in the direction of their momentum by their energy relaxation time τ_ε, and e and μ are the charge and chemical potential of a conduction electrons.

2. THERMOELECTRIC EFFECTS IN A MAGNETIC FIELD ORTHOGONAL TO THE LAYERS

The FS of most layered conductors has many sheets and consists of topologically different elements in the form of cylinders and planes weakly corrugated along axis $p_z = \mathbf{p}\cdot\mathbf{n}$, where **n** is the normal to the layers. As in [3], we assume that the FS in each unit cell of the momentum space consists of a cylinder and two quasi-planar surfaces adjoining it; the normal to the quasi-planar FS sheets will be referred to as the p_x axis (Fig. 1). Such a topological structure is typical of a large family of organic conductors based on tetrathiafulvalene. The weakly corrugated cylinder is equidistant from the quasi-planar FS sheets, and its sections by plane $p_B = \mathbf{p}\cdot\mathbf{B}/B = \text{const}$ are symmetric in a magnetic field orthogonal to the layers.

In the immediate vicinity to the Lifshits topological transition, under the action of an external perturbation, minimal distance Δ_p between individual FS sheets (cavities) turns out to be so small that a conduction electron can pass from one FS sheet to the other because of magnetic breakdown with probability

$$w(p_B) = \exp(-2cS_p/eB\hbar), \qquad (2)$$

where $S_p \approx \Delta_p^2$ is the over-barrier area which the electron must cross during tunneling from one FS sheet to the other [10].

Among chaotic walks of the electron over different FS cavities, the strictly periodic motion of the charge turns out to be most probable, when, during each probable magnetic breakdown, the electron necessarily passes to the adjacent FS cavity if during its mean free time it manages to visit all segments of its magnetic-breakdown trajectory many times. The mathematical expectation of such a finite motion of the electron is close to unity if the following condition is satisfied [11, 12]:

$$w(p_B) \gg \gamma \equiv 1/\omega_c\tau. \qquad (3)$$

The quantized electron energy levels can easily be determined using the Bohr–Sommerfeld rule of quantization of area $S(\varepsilon, p_B)$ in the momentum space, which is circumvented by the electron during its motion in a closed orbit with conservation of its energy $\varepsilon(\mathbf{p})$ and projection p_B of the momentum onto the magnetic field direction,

$$S(\varepsilon, p_B) = 2\pi\hbar\frac{eB}{c}\left(n+\frac{1}{2}\right),\quad n = 0, 1, 2, 3, \ldots, \tag{4}$$

where c is the velocity of light.

In a quantizing magnetic field, peculiarities appear in the density of states of charge carriers,

$$\nu(\varepsilon) = \int\frac{2dp_B\partial S(\varepsilon, p_B)}{(2\pi\hbar)^3\partial\varepsilon}, \tag{5}$$

for magnetic field values B_n for which $\partial S(\varepsilon, p_B)/\partial p_B$ vanishes. This can easily be verified by expanding the left-hand side of Eq. (4) into a power series in $\delta p_B = p_B - p_B^0$ near the electron orbit with extremal area $S_{\text{ext}} = S(\varepsilon, p_B^0)$:

$$S(\varepsilon, p_B^0) + \frac{\partial^2 S(\varepsilon, p_B^0)}{(\partial p_B^0)^2}\frac{\delta p_B^2}{2} = 2\pi\hbar\frac{eB}{c}\left(n+\frac{1}{2}\right). \tag{6}$$

It can easily be noted that the root singularities of function $\nu(\varepsilon)$, which are associated with the extremum S_{ext} of the area of the closed orbit in the momentum space, are formed by a small number of conduction electrons, $\delta p_B \le \sqrt{2\pi\hbar eB/c\eta}$, or by their relative fraction

$$a\frac{\delta p_B}{\hbar} \le \sqrt{\frac{\hbar\omega_c}{\mu\eta}} \ll 1, \tag{7}$$

where a is the distance between the layers and $\eta = \partial^2 S(\varepsilon, p_B^0)/(\partial p_B^0)^2$ is the quasi-two-dimensionality parameter of the electron energy spectrum, which coincides (to within a numerical factor of the order of unity) with the ratio of maximal velocity $\overline{v}_z$ of the charge carrier drift along normal $\mathbf{n}$ to the layers to characteristic Fermi velocity v_F of electrons moving over the layers. These singularities are repeated periodically upon a change in the reciprocal magnetic field in accordance with relation (4) with period

$$\Delta\left(\frac{1}{B}\right) = \frac{2\pi\hbar e}{cS_{\text{ext}}}, \tag{8}$$

which just leads to a periodic variation of thermodynamic and kinetic characteristics of degenerate conductors as a function of $1/B$.

If not only the magnetic field, but also the temperature gradient is orthogonal to the layers, electric field E_z in the main approximation in the small quasi-two-dimensionality parameter of the conductor ($\eta \ll 1$),

$$E_z = \frac{\pi^2 T}{3e\sigma_{zz}(\mu)}\frac{\partial\sigma_{zz}(\mu)}{\partial\mu}\frac{\partial T}{\partial z}, \tag{9}$$

is an asymptotic function of only the conductivity tensor component

$$\sigma_{zz} = -\sum_{k=-\infty}^{\infty}\int d\varepsilon\frac{\partial f_0(\varepsilon)}{\partial\varepsilon}\int\frac{2dp_B}{(2\pi\hbar)^3}(e\overline{v}_z)^2 \times\tau_B\frac{\partial S(\varepsilon, p_B)}{\partial\varepsilon}\exp\left(\frac{ikcS(\varepsilon, p_B)}{\hbar eB}\right). \tag{10}$$

Here, $f_0 = \{1 + \exp(\varepsilon - \mu)/T\}^{-1}$ is the equilibrium Fermi distribution function for conduction electrons. Using the Poisson relation [13]

$$\sum_{n=0}^{\infty}\varphi_n = \sum_{k=-\infty}^{\infty}\int_0^{\infty}dn\varphi(n)\exp(2\pi ikn),$$

we have replaced summation over n by integration with respect to n in summing over all electron states defined by variables p_B and n and then, using relation

$$dn = \frac{c}{2\pi\hbar eB}\frac{\partial S(\varepsilon, p_B)}{\partial\varepsilon}d\varepsilon \tag{11}$$

by integration with respect to ε. Electron mean free time $\tau_B = \tau(1 + \nu_{\text{osc}})$ contains the following correction oscillating with the magnetic field:

$$\nu_{\text{osc}} = \sum_{q=1}^{\infty}\xi_q\int d\varepsilon\frac{\partial f_0(\varepsilon)}{\partial\varepsilon}\int\frac{adp_B'}{\hbar}\cos\frac{qcS(\varepsilon, p_B')}{eB\hbar}, \tag{12}$$

which is connected with quantum oscillations of the electron scattering amplitude [14]. Numerical factors ξ_q of the order of unity in formula (12) depend on the specific form of the scattering amplitude of charge carriers and their dispersion relation. In the case of an isotropic electron energy spectrum, these factors were calculated by many authors using various methods [15–18]. However, the knowledge of these factors is immaterial for solving the inverse problem of reconstructing the energy spectrum of conduction electrons from the results of experimental investigation of electron phenomena in a quantizing magnetic field. It is sufficient to know only extremal closed planar sections of the FS.

The oscillatory dependence of the conductivity tensor components for quasi-two-dimensional conductors on the inverse quantizing magnetic field is much smaller than the smoothly varying part of these components for $\hbar\omega_c/\mu \ll \eta \ll 1$ (approximately by a factor of $\sqrt{\mu\eta/\hbar\omega_c}$). Differentiation of the rapidly oscillating exponential function $\exp[ikcS(\mu, p_B)/\hbar eB]$

in expression (10) with respect to μ leads to multiplication of this exponential by quantity

$$\frac{ikc}{\hbar eB}\frac{\partial S}{\partial \mu} = \frac{2\pi ik}{\hbar\omega_c},$$

and the term with $k = 0$ acquires a much smaller factor on the order of $1/\mu$. As a result, quantum oscillations of the thermoelectric field are giant by nature, and oscillating part $\mathbf{E}_{\rm osc}(B)$ of thermoelectric field is much larger than the nonoscillating part $\mathbf{E}^{\rm mon}(B)$ that coincides with the thermoelectric field for $\hbar\omega_c \ll T$.

The presence of two sharp functions in formulas (10) and (12) makes it possible to easily integrate them with respect to ε, p'_B, and p_B and obtain (for $2\pi^2 T \ll \hbar\omega_c$ and $\hbar\omega_c/\mu \ll \eta \ll 1$) the following compact asymptotic expression for the thermoelectric field:

$$E_z = E_z^{\rm mon}\left[1 + \xi\sqrt{\frac{\mu}{\hbar\omega_c}}\sin\left(\frac{cS_{\rm ext}}{eB\hbar} + \frac{s\pi}{4}\right)\right], \tag{13}$$

where $s = {\rm sgn}[\partial^2 S(\varepsilon, p_B)/\partial p_B^2]$ and ξ is a dimensionless factor on the order of unity, which depends on the dispersion relation for charge carriers; the smoothly varying part of the thermoelectric field in layered conductors has form

$$E_z^{\rm mon} = \frac{\pi^2 T}{3e}\frac{\partial}{\partial z}\ln\int dp_B[\overline{v_z(\mu, p_B)}]^2\tau \tag{14}$$

and has the same order of magnitude as $T/e\mu$. Here, we have written only the first harmonics of oscillating functions with $s = \pm 1$ (i.e., the main contributions to quantum oscillations of electrons in neighborhoods with maximal and minimal sections of FS by plane $p_B = {\rm const}$).

We assume that the maximal section of the FS cylindrical part is much closer to FS quasi-planar sheets (as shown in Fig. 1a) than the minimal section. In this case, electrons near minimal section $S_c^{\rm min}$ do not leave the FS cylindrical part and form conventional Shubnikov–de Haas oscillations with frequency $\nu = cS_c^{\rm min}/e\hbar$ proportional to $S_c^{\rm min}$ (green region in Fig. 1), while electrons on the FS cylindrical part near maximal area $S_c^{\rm max}$ of its section participate together with charge carriers on the FS quasi-planar sheets in the formation of magnetic-breakdown thermoelectric field oscillations. The area of the magnetic-breakdown electron trajectory (red region in Fig. 1) consists of four pieces between magnetic-breakdown contacts $(\pm D/2, 0, 0)$ and $(\pm D/2, P_2, 0)$, where D is the diameter of the maximal section of the FS cylindrical part along the p_x axis, $P_2 = 2\pi\hbar/a_2$ is the period of the unit cell in the momentum space along the p_y axis, and a_2 is the period of the crystal in the conventional space along the y axis.

The origin in the p_yp_z plane coincides with one of the magnetic-breakdown contacts. The energy levels of conduction electrons moving in closed magnetic-breakdown orbits can be determined using Eq. (4), substituting the total area circumvented by an electron moving over two adjacent cylinders and over both FS quasi-planar sheets:

$$S_{mb} = S_c^{\rm max} + S_{pl}, \tag{15}$$

where

$$S_{pl}(\varepsilon, p_B^0) = \int_{-D/2}^{D/2} dp_x[p_y^{(2)}(p_x) - p_y^{(1)}(p_x)] \tag{16}$$

is the area between the two FS quasi-planar sheets per unit cell in the momentum space.

In the case of nonuniform heating of the sample along the layers, nonoscillating part $E_y^{\rm mon}(B)$ of the thermoelectric field increases linearly with the magnetic field when the temperature gradient is directed along the x axis (see formula (55) in [4]). In all remaining cases, the Nernst–Ettingshausen field, as well as the thermoelectric field along the temperature gradient, oscillates against the background of a weak constant field on the order of $T/e\mu$.

3. THERMOELECTRIC EFFECTS IN A MAGNETIC FIELD TILTED TO THE LAYERS

In magnetic field $\mathbf{B} = (B\sin\vartheta\cos\varphi, B\sin\vartheta\sin\varphi, B\cos\vartheta)$ tilted from the normal to the layers, section $S_c(\varepsilon, p_B)$ of the FS cylindrical cavity is extended along the p_z axis. The electron time of flight $T_1(\vartheta) = T_1(0)/\cos\vartheta$ between the FS quasi-planar sheets increases with the tilt angle of the magnetic field to the layers, and the electron turning points on the closed orbit along the p_x axis are at different distances from the FS planar sheets. For small angles ϑ, this difference increases and may become equal to corrugation $\eta\hbar/a$ of the FS cylindrical part. If $\exp(-c\hbar\eta^2/a^2eB)$ in this case is smaller than γ, an electron moving in the magnetic field lying in the xz plane is unable to close the magnetic-breakdown orbit, and its motion over FS planar sheets is infinite. In this case, the electron mainly moves in an open trajectory, sometimes completing a turn on the closed section of the FS cylindrical part, and there are no quantum magnetic-breakdown oscillations of kinetic coefficients.

In this case, quantum oscillations of kinetic coefficients are formed only by electrons on the FS cylindrical part. For $\tan\vartheta \gg 1$, classical angular oscillations also come into play; however, their amplitude is much smaller than the amplitudes of quantum oscillations that suppress classical thermoelectric field oscillations. However, there exists a countable set of magnetic field orientations, for which magnetic-breakdown quantum oscillations of the same intensity as in the magnetic field orthogonal to the layers appear again. This is possible when the four magnetic-breakdown contacts lie in the same plane.

This can occur when an electron moving from the FS cylindrical part to a quasi-planar FS sheet and continuing its motion encounters a magnetic-breakdown contact the same as for $\vartheta = 0$ on the adjacent FS cylinder or at least at the next cylinder. Passing to this neighboring cylinder, the electron must complete an integer number of periods P_3 along the p_z axis during its motion and must then return to the previous FS cylinder and close the magnetic-breakdown trajectory. In view of the periodic dependence of the energy of the electron on its quasi-momentum, the magnetic breakdown probability turns out to be the same and completely identical to that in the magnetic field orthogonal to the layers.

A conduction electron that has begun its motion, say, at contact $p_1 = (D_p/2, 0, 0)$ moves in an open orbit to contact $p_2 = (D_p/2, mP_2, nP_3)$ after shifting by an integer number of unit cell periods P_2 and P_3 and then rises over the FS cylindrical part along the p_z axis by an integer number N of periods P_3 and shifts to the FS opposite sheet to contacts $p_3 = (-D_p/2, mP_2, (n + N)P_3)$. Having shifted down in an open trajectory by n periods P_3, the electron approaches contact $p_4 = (-D_p/2, 0, NP_3)$ and then closes the orbit at contact $p_1 = (D_p/2, 0, 0)$.

The magnetic field orientation orthogonal to such a closed orbit is determined by the conditions of orthogonality of magnetic field vector **B** to vectors connecting contact $p_1 = (D_p/2, 0, 0)$ with the remaining aforementioned magnetic-breakdown contacts. This orthogonality conditions implies that

$$\cos\varphi\tan\vartheta = \frac{2\pi\hbar N}{aD_p}, \quad \sin\varphi\tan\vartheta = \frac{na_2}{ma}. \qquad (17)$$

It is convenient to transform these relations to

$$\tan\vartheta = \sqrt{\left(\frac{2\pi\hbar N}{aD}\right)^2 + \left(\frac{na_2}{ma}\right)^2}, \quad \tan\varphi = \frac{na_2 D}{2\pi\hbar Nm}. \qquad (18)$$

For the magnetic field orientation satisfying this condition, the amplitude of magnetic-breakdown oscillations of kinetic coefficients is of the same order of magnitude as in the magnetic field normal to the layers because the probability that the electrons passes to the other FS sheet as a result of magnetic breakdown in both directions of the magnetic field (normal and tilted to the layers) is determined only by width Δ_p of the gap between the FS sheets. In the case of slight violation of condition (18), the magnetic-breakdown oscillation amplitude begins to decay and assumes the minimal value in the collisionless limit ($\tau = \infty$) for magnetic field tilt angle ϑ_c to the layers that satisfy condition (18) if we substitute half-integer number $N + 1/2$ for N or $n + 1/2$ for n.

Magnetic-breakdown quantum oscillations of resistivity of organic conductors was detected for the first time in organic complex κ ((BEDT—TTF)$_2$Cu(NCS)$_2$) then in other layers of tetrathiofulvene also at different laboratories [19, 20]. The resistivities of this family of organic conductors across and along the layers differ by three orders of magnitude. In such complexes of organic conductors, giant oscillations of the thermoelectric field with a change in the reciprocal value of the quantizing magnetic field by about 20–30 T and their periodic disappearance as a function of $\tan\vartheta$ can easily be detected.

The strong magnetic field condition is determined by the number of turns of an electron in a closed orbit during its mean free time, which is inversely proportional to the area of the electron orbit. More reliable information is contained in the results of measurement in the case of large orbits with small values of N and n, including $N = 0$ or $n = 0$. For $n = 0$ (i.e., for magnetic field rotation in the xz plane), sharp thermoelectric field magnetic-breakdown peaks appear with frequency

$$\nu = \frac{c(S_c^{\max} + S_{pl})}{e\hbar\cos\vartheta_N}$$

for the magnetic field tilt angle satisfying condition

$$\tan\vartheta_N = 2\pi\hbar N/aD, \quad N = 1, 2, 3, \qquad (19)$$

and with frequency

$$\nu_m = \frac{c(S_c^{\max} + mS_{pl})}{e\hbar\cos\vartheta_m}$$

for $N = 0$, $n = 1$, when the magnetic field lies in the yz plane, and its tilt angle to the layers satisfies condition

$$\tan\vartheta_m = a_2/ma. \qquad (20)$$

Experimental investigation of these oscillatory effects will make it possible to determine the area of the maximal section of the FS cylindrical part and its diameter D along the p_x axis. Comparing oscillation frequency ν_2 for $N = 0$, $n = 1$, $m = 2$ with the magnetic-breakdown orbit depicted in Fig. 1b with the frequency of magnetic-breakdown oscillations in a magnetic field orthogonal to the layers,

$$\nu_0 = \frac{c(S_c^{\max} + S_{pl})}{e\hbar},$$

i.e., for $N = 0$, $n = 0$, we can determine area S_{pl} between quasi-planar FS sheets per unit cell of the momentum space,

$$S_{pl} = \frac{e\hbar}{c}(\nu_2\cos\vartheta_2 - \nu_0). \qquad (21)$$

In the case of nonuniform heating of the conductor along the layers, Shubnikov–de Haas magnetic-breakdown oscillations of thermoelectric field,

$$E_z = \frac{\pi^2 T}{3e}\rho_{za}\frac{\partial\sigma_{\alpha\beta}(\mu)}{\partial\mu}\frac{\partial T}{\partial r_\beta},$$
$$E_a = \frac{\pi^2 T}{3e}\rho_{a\delta}\frac{\partial\sigma_{\delta\beta}(\mu)}{\partial\mu}\frac{\partial T}{\partial r_\beta}, \quad (\alpha,\beta,\delta) = (x,y), \tag{22}$$

occur against the background that increases monotonically with the magnetic field like in the case of the magnetic field normal to the layers. However, for solving the inverse problem of reconstruction of the FS, it is sufficient to know only the periods of quantum and angular oscillations of the thermoelectric field.

4. CONCLUSIONS

The giant nature of thermoelectric field quantum oscillations makes it possible to determine specific characteristics of the electron energy spectrum of layered conductors with a high degree of accuracy (in particular, to measure the areas and diameters of electron orbits on the FS for various orientations of the strong magnetic field during a slow approach to a topological transition using continuous and quite controllable variation of pressure). With increasing temperature, at transition to the temperature range of liquid hydrogen, quantum oscillations of kinetic coefficients begin to decay exponentially and cannot prevent the observation of classical angular oscillations, which also contain important information about the energy spectrum of charge carriers.

As a result of complex investigation of thermoelectric phenomena at various temperatures that are much lower than the Debye temperature, it is quite possible to solve the inverse problem of reconstruction (from experimental data) of the FS, which is the basic characteristic of the electron energy spectrum of layered conductors, and to get information on relaxation properties of conduction electrons.

REFERENCES

1. I. M. Lifshits, Sov. Phys. JETP **9**, 1107 (1960).
2. N. V. Zavaritsky, in *Selected Works of I. M. Lifshits. Electronic Theory of Metals. Polymers and Biopolymers* (Nauka, Moscow, 1994), p. 432 [in Russian].
3. V. G. Peschanskii, O. Galbova, and R. Hasan, J. Exp. Theor. Phys. **123**, 1060 (2016).
4. V. G. Peschansky, Low Temp. Phys. **43**, 303 (2017).
5. L. D. Landau, Z. Phys. **64**, 629 (1930).
6. W. J. de Haas and P. M. van Alphen, Leiden Commun. **212a**, 215 (1930).
7. L. W. Shubnikov and W. J. de Haas, Leiden Commun. **210a**, 3 (1930).
8. D. Shoenberg, Proc. R. Soc. London, Ser. A **170**, 363 (1939).
9. I. M. Lifshits and A. M. Kosevich, Sov. Phys. JETP **2**, 636 (1955).
10. G. A. Gamov, Usp. Fiz. Nauk **10**, 531 (1930).
11. O. Galbova, O. V. Kirichenko, and V. G. Peschansky, Low Temp. Phys. **39**, 780 (2013).
12. V. G. Peschansky and D. I. Stepanenko, J. Exp. Theor. Phys. **123**, 156 (2016).
13. R. Courant, *Methods of Mathematical Physics* (Wiley, New York, 1989).
14. B. Davydov and I. Pomeranchuk, J. Phys. USSR **2**, 147 (1940).
15. G. E. Zil'berman, Sov. Phys. JETP **2**, 650 (1955).
16. R. Kubo, H. Nasegava, and N. Hashitsume, J. Phys. Jpn. **14**, 56 (1959).
17. Yu. A. Bychkov, Sov. Phys. JETP **12**, 971 (1960).
18. A. M. Kosevich and V. V. Andreev, Sov. Phys. JETP **12**, 608 (1960).
19. K. Andres, C.-P. Heidmann, H. Müller, S. Himmellsbach, W. Biberacher, Ch. Probst, and W. Joss, Synth. Met. **41–43**, 1893 (1991).
20. C.-P. Heidmann, H. Müller, W. Biberacher, K. Neumaier, Ch. Probst, K. Andres, A. G. M. Jansen, and W. Joss, Synth. Met. **41–43**, 2029 (1991).

Translated by N. Wadhwa