ЖУРНАЛ ФІЗИЧНИХ ДОСЛІДЖЕНЬ
т. 13, № 4 (2009) 4704(5 с.)

JOURNAL OF PHYSICAL STUDIES
v. 13, No. 4 (2009) 4704(5 p.)

ELECTRON TRANSPORT IN STRONGLY ANISOTROPIC STRUCTURES IN A MAGNETIC FIELD

O. V. Kirichenko[1], V. G. Peschansky[1], O. Galbova[2]

[1] *B. Verkin Institute for Low Temperature Physics and Engineering, National Academy of Sciences of Ukraine, Lenin Ave. 47, 61103, Kharkiv, Ukraine*

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

(Received October 16, 2009)

A linear response of the electron system of a layered conductor in the presence of a temperature gradient is investigated theoretically. We have found the dependencies of the thermo-emf on the temperature and on an external magnetic field, the experimental study of which allows to determine the structure of charge carriers energy spectrum as well as to examine different relaxation mechanisms in the system of conduction electrons.

Key words: layered conductor, thermoelectric field, magnetic field.

PACS number(s): 72.15.Jf

Investigations of the thermoelectric effect in layered conductors placed in a strong magnetic field $\mathbf{B}$ allow to study in detail their electron energy spectrum [1, 2].

Within the framework of the linear theory the electric current in a conductor

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

is related to an electric field $\mathbf{E}$ and a temperature gradient ∇T by the conductivity tensor σ_{ik} and the thermoelectric tensor α_{ik}.

In the case of nonuniform temperature distribution in a specimen, even in the absence of current-conducting contacts ($\mathbf{j} = 0$), the thermoelectric field

$$E_i = \rho_{il}\ \alpha_{lj} \frac{\partial T}{\partial x_j}, \tag{2}$$

unavoidably appears. Here ρ_{ij} is the resistivity tensor which is inverse to the conductivity tensor σ_{ij}.

Some important information about relaxation processes in the electron system may be extracted from the T-dependence of the thermoelectric field. The reason is that the kinetic coefficients

$$\sigma_{ik} = \frac{2e^3 B}{c(2\pi\hbar)^3} \int \frac{\partial f_0(\varepsilon)}{\partial \varepsilon} d\varepsilon \int dp_B \times \int_{-\infty}^{0} \exp(t/\tau_p) dt \int_0^{T_B} dt' v_i(t') v_k(t'+t) \tag{3}$$

and

$$\alpha_{ik} = \frac{2e^2 B}{c(2\pi\hbar)^3} \int \frac{\partial f_0(\varepsilon)}{\partial \varepsilon}\ \frac{\varepsilon - \mu}{T} d\varepsilon \int dp_B \times \int_{-\infty}^{0} \exp(t/\tau_\varepsilon) dt \int_0^{T_B} dt' v_i(t') v_k(t'+t) \tag{4}$$

describe different relaxation processes in the charge carriers system. The components of the conductivity tensor are connected to the momentum relaxation of electrons which is characterized by the time τ_p, and the components of the tensor α_{ij} depend on the energy relaxation the time τ_ε. Here $e, \mathbf{v}, p_B, \varepsilon$ are the charge, velocity, momentum projection onto $\mathbf{B}$-direction, and energy of a conduction electron, respectively; $f_0(\varepsilon)$ is the equilibrium Fermi distribution function, μ is the chemical potential of the electron system, c is the velocity of light, $\hbar$ is the Planck constant, t is the time of charge motion in a magnetic field according to the equation

$$\frac{\partial \mathbf{p}}{\partial t} = \frac{e}{c} \mathbf{v} \times \mathbf{B}.$$

In case of a periodic motion of an electron in the magnetic field the magnitude T_B is the period of motion. If an electron executes an unperiodic motion along an open trajectory, T_B is the time that characterizes its displacement per period of the inverse lattice.

In the absence of the magnetic field the thermoelectric field

$$E_i = \frac{\pi^2}{3e} \left(\frac{T}{\mu}\right) Q_{ij}\ \frac{\tau_\varepsilon}{\tau_p}\ \frac{\partial T}{\partial x_J} \tag{5}$$

is proportional to the ratio of the relaxation times $\tau_\varepsilon / \tau_p$. Here the dimensionless tensor Q_{ij} does not depend on temperature.

We consider the thermoelectric effect in a quasi-two-dimensional conductor at temperature that is much below the Debye temperature T_{D}, when the T-dependencies of the relaxation times τ_ε and τ_p are essentially different. When the temperature is close to zero, the relaxation in electron system of a degenerated conductor is realized mainly by the charge carriers scattering at impurity centers and other crystal defects. This is the case when the relaxation times τ_p and τ_ε are of the same order of magnitude. With the temperature increasing an extra mechanism of relaxation connected to electron scattering by crystal lattice vibrations sets off. According to Matthiessen's rule each of the scattering mechanisms makes an additive contribution into the relaxation process, so

$$\frac{1}{\tau} = \frac{1}{\tau^{(\mathrm{im})}} + \frac{1}{\tau^{(\mathrm{eph})}}, \tag{6}$$

where $1/\tau^{(\mathrm{eph})}$ is the frequency of charge carriers collisions with phonons. At $T \ll T_\mathrm{D}$ because of the small-angle scattering of electrons by phonons, a number of collisions with phonons, necessary for the electron momentum relaxation, is much greater than that for the energy relaxation. As a result, τ_ε decreases more rapidly with temperature than τ_p . The T-dependence of τ_ε has the form

$$\tau_\varepsilon = \left\{ \frac{1}{\tau^{(\mathrm{im})}} + \frac{1}{\tau_0}\, \tilde{T}^n \right\}^{-1}, \tag{7}$$

where $\tilde{T} = \frac{T}{T_\mathrm{D}}$, τ_0 is the time characterizing the energy relaxation of electrons due to their collisions with phonons at the Debye temperature. The power n is determined by the dimensionality of the system, n equals 3 for 3D metals and n equals 2 in the case of 2D conductors [3].

At sufficiently low temperatures, when the momentum relaxation is realized mainly due to the charge carriers collisions with impurities (i. e. τ_p coincides with $\tau^{(\mathrm{im})}$) and account of the interaction with phonons affects essentially the energy relaxation, the T-dependence of the thermoelectric field at $\mathbf{B} = 0$ can be presented as

$$E_i = \frac{\pi^2 T_\mathrm{D}}{3e\mu} Q_{ik} f(\tilde{T})\, \frac{\partial T}{\partial x_k},$$
$$f(\tilde{T}) = \frac{\tilde{T}}{1 + (\tau^{(\mathrm{im})}/\tau_0)\tilde{T}^n}. \tag{8}$$

At $\tilde{T} \ll (\tau_0/\tau^{(\mathrm{im})})^{1/n}$ the field E_i is proportional to T, but with the temperature increasing the frequency of charge carriers collisions with phonons $1/\tau_\varepsilon^{(\mathrm{eph})} = 1/(\tau_0)\tilde{T}^n$ becomes comparable to the frequency of their collisions with impurities $1/\tau^{(\mathrm{im})}$ and the growth with temperature gives place to decreasing as $T^{(1-n)}$. The maximum is attained at

$$\tau_\varepsilon^{(\mathrm{eph})} = (n-1)\tau^{(\mathrm{im})}.$$

In layered conducting structures with the Q2D energy spectrum at temperature that is much less than the overlap integral $t_\perp$ for the wave functions of electrons belonging to neighboring layers, the T-dependence of $\tau_\varepsilon^{(\mathrm{eph})}$ is basically the same as in 3D metals. In layered organic conductors the overlap integral $t_\perp$ is of the order of the Debye temperature. Thus in the whole of the range where T is less than the Debye temperature, $1/\tau_\varepsilon^{(\mathrm{eph})}$ is proportional to T^3, and in accordance with the Bloch law, τ_p is of the form

$$\frac{1}{\tau_p} = \frac{1}{\tau^{(\mathrm{im})}} + \frac{1}{\tau_0}\, \tilde{T}^5. \tag{9}$$

In certain organic conductors there are several groups of charge carriers that are responsible for the electron transport. The Fermi surface in such conductors may consist of topologically different elements: weekly corrugated cylinders and planes [4, 5].

An external magnetic field affects differently the motion of charge carriers whose states belong to the weakly corrugated cylinder and to the plane sheet of the Fermi surface. Just for this reason the presence of such plane sheets of the Fermi surface may be revealed most easy in a conductor placed in a magnetic field.

As an example, consider the conductor with two groups of charge carriers, in a magnetic field $\mathbf{B} = (0, B\sin\theta, B\cos\theta)$. It is assumed that the preferred directions of the mean velocities $\pm\mathbf{v}_1$ of electrons belonging to the plane sheets, are determined by the angle φ, so $v_{1x} = \pm v_1\cos\varphi$, $v_{1y} = \pm v_1 \sin\varphi$, and the dispersion law for the other charge carriers group is of the form

$$\varepsilon(\mathbf{p}) = \frac{p_x^2 + p_y^2}{2m} - 2t_\perp \cos\frac{ap_z}{\hbar}, \tag{10}$$

where $m = \mathrm{const}$, a is the separation between the layers, $t_\perp$ is much less than the Fermi energy ε_F.

We shall assume that the angle θ is not near $\pi/2$, so all orbits of electrons with the Q2D dispersion law are closed. At $\cos\theta \gg mc/eB\tau$ the condition of the strong magnetic field $(T_B/\tau) \ll 1$ is fulfilled automatically.

The components of the kinetic coefficients for the conductor σ_{ik} and α_{ik} are the sums of the contributions of each group of charge carriers, and, in particular, it is as easy to make sure that

$$\sigma_{ik} = \begin{pmatrix} \gamma^2\sigma_2 - \gamma^2\sigma_{zz}\tan^2\theta + \sigma_1\cos^2\varphi & \gamma\sigma_2 - \gamma\sigma_{zz}\tan^2\theta + \sigma_1\cos\varphi\sin\varphi & -\gamma\sigma_{zz}\tan\theta \\ -\gamma\sigma_2 + \gamma\sigma_{zz}\tan^2\theta + \sigma_1\cos\varphi\sin\varphi & \gamma^2\sigma_2 + \sigma_{zz}\tan^2\theta + \sigma_1\sin^2\varphi & \sigma_{zz}\tan\theta \\ \gamma\sigma_{zz}\tan\theta & \sigma_{zz}\tan\theta & \sigma_{zz} \end{pmatrix}. \tag{11}$$

Here σ_1 and $\sigma_2 = (e^2\varepsilon_\mathrm{F}\tau_p/\pi\hbar^2 a)$ are the contributions to the conductivity along the layers at $\mathbf{B} = 0$ made by charge carriers belonging to the plane sheets of the Fermi surface and to the corrugated cylinder, respectively; $\gamma = mc/(eB\tau_p\cos\theta) = 1/(\omega_c\tau_p) \ll 1$.

Account of the contribution to σ_{zz} of charge carriers with Q1D energy spectrum does not affect the T-dependence of the thermoelectric field. The component

$$\sigma_{zz} = \frac{2ae^2 m\tau_p t_\perp^2 \cos\theta}{\pi\hbar^4}\; J_0^2\left(\frac{ap_\mathrm{F}}{\hbar}\tan\theta\right) \equiv S_z\tau_p, \tag{12}$$

in the main approximation in the quasi-two-dimensionality parameter $t_\perp/\varepsilon_{\rm F} = \eta \ll 1$ is quadratic in η, and at $1 \ll \tan\theta \ll eB\tau_p/mc$ this component oscillates as θ varies. Here $J_n(x)$ is the Bessel function, $p_{\rm F}$ is the Fermi momentum.

At $\gamma \ll 1$ the asymptotic expression for the resistivity tensor is

$$\rho_{ik} = \rho_0 \begin{pmatrix} 1 + q\dfrac{\sin^2\varphi}{\gamma^2} & -\dfrac{1}{\gamma} - q\dfrac{\sin 2\varphi}{2\gamma^2} & \left(\dfrac{1}{\gamma} + q\dfrac{\sin 2\varphi}{2\gamma^2}\right)\tan\theta \\ \dfrac{1}{\gamma} - q\dfrac{\sin 2\varphi}{2\gamma^2} & 1 + q\dfrac{\cos^2\varphi}{\gamma^2} & -q\dfrac{\cos^2\varphi}{\gamma^2}\tan\theta \\ \left(-\dfrac{1}{\gamma} + q\dfrac{\sin 2\varphi}{2\gamma^2}\right)\tan\theta & -q\dfrac{\cos^2\varphi}{\gamma^2}\tan\theta & \dfrac{1}{\sigma_{zz}^{(2)}\rho_0} + q\dfrac{\cos^2\varphi}{\gamma^2}\tan^2\theta \end{pmatrix}, \tag{13}$$

where $\rho_0 = 1/(\sigma_1 + \sigma_2)$, $q = \sigma_1/\sigma_2$.

The thermoelectric effect depends essentially on the presence of the plane sheet of the Fermi surface, and the distinct components of the field

$$E_i = \frac{\pi^2 T_{\rm D}}{3e\mu} P_{ik}(T)\frac{\partial T}{\partial x_k} \tag{14}$$

behave in different ways.

In the case of a single charge carriers group with the Q2D dispersion law, in the main approximation in the small parameters γ and η the diagonal components

$$P_{xx} = P_{yy} = \tilde{T} \tag{15}$$

exceed significantly the components

$$P_{xy} = -P_{yx} = \tilde{T}\left(\frac{1}{\omega_c\tau_p} - \frac{1-b}{\omega_c\tau_\varepsilon}\right), \tag{16}$$

so the field in the layers plane is directed mainly along ∇T and grows linearly with T, when a conductor is heated along the layers. The components

$$P_{zx} = -\tilde{T}\left(\frac{1-b}{\omega_c\tau_\varepsilon} - \frac{d}{\omega_c\tau_p}\right)\tan\theta, \tag{17}$$

and

$$P_{zy} = \tilde{T}\left(-1 + \frac{(d+b)\tau_\varepsilon}{\tau_p}\right)\tan\theta \tag{18}$$

differ from zero if only the magnetic field deviates from the normal to the layers,

$$P_{zz} = \tilde{T}(d+b), \tag{19}$$

and the magnitudes P_{xz}, P_{yz} are proportional to η^2. Here $b = (\mu/\tau_\varepsilon)(\partial\tau_\varepsilon/\partial\mu) \simeq 1$, and the coefficient $d = (\mu/S_z)(\partial S_z/\partial\mu)$ is of the form

$$d = -\frac{2\mu ma\tan\theta}{p_{\rm F}\hbar} \times \frac{J_1\left(\dfrac{ap_{\rm F}}{\hbar}\tan\theta\right)J_0\left(\dfrac{ap_{\rm F}}{\hbar}\tan\theta\right)}{J_0^2\left(\dfrac{ap_{\rm F}}{\hbar}\tan\theta\right) + \phi_1\gamma^2 + \phi_2\eta^2}, \tag{20}$$

where ϕ_1 and ϕ_2 are the magnitudes of the order of unit which allow for corrections to S_z, omitted in the main approximation in the small parameters η and γ. As a result, the dependence of the components P_{zk} upon the angle θ takes form of giant oscillations. The dependence of the nondiagonal components of the tensor P_{ik} on the temperature is determined by the T-dependence of the relaxation times, and taking into consideration (7) and (9),we have

$$P_{xy} = \frac{\tilde{T}}{\omega_c\tau^{(\rm im)}}\left[b - (1-b)\frac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^3 + \frac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^5\right], \tag{21}$$

$$P_{zx} = \frac{\tilde{T}}{\omega_c\tau^{(\rm im)}}\left[d + b - 1 - (1-b)\frac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^3 + d\frac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^5\right]\tan\theta, \tag{22}$$

$$P_{zy} = \tilde{T}\left[-1 + (d+b)\frac{1 + \dfrac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^5}{1 + \dfrac{\tau^{(\rm im)}}{\tau_0}\tilde{T}^3}\right]\tan\theta. \tag{23}$$

In the presence of the Q1D charge carriers group the components of the tensor P_{ik} contain the terms, that grow linearly with B increasing. These terms vanish at certain values of the angle φ, in the remaining cases they are dominant and determine the thermoelectric effect in the strong magnetic field ($\omega_c\tau \gg 1$). When ∇T is oriented in the plane of the layers the linear growth with B of the components

$$P_{xx} = \rho_0\tilde{T}\left[\sigma_2 + \sigma_1\left((1-b)\frac{\tau_p}{\tau_\varepsilon}\sin^2\varphi + a\frac{\tau_\varepsilon}{\tau_p}\cos^2\varphi + \omega_c(\tau_p - a\tau_\varepsilon)\sin\varphi\cos\varphi\right)\right], \tag{24}$$

$$P_{yy} = \rho_0\tilde{T}\left[\sigma_2 + \sigma_1\left((1-b)\frac{\tau_p}{\tau_\varepsilon}\cos^2\varphi + a\frac{\tau_\varepsilon}{\tau_p}\sin^2\varphi - \omega_c(\tau_p - a\tau_\varepsilon)\sin\varphi\cos\varphi\right)\right] \tag{25}$$

takes place when both projections (v_{1x} and v_{1y}) of the electron velocity at the plane sheet of the Fermi surface vanish simultaneously. This is the case when the T-dependence of the components P_{xx}, P_{yy} in the main approximation in the parameter γ has the form

$$P_{xx} = -P_{yy} = \rho_0\sigma_1\omega\tau^{(\mathrm{im})}\tilde{T}\left(\frac{1}{1+\dfrac{\tau^{(\mathrm{im})}}{\tau_0}\tilde{T}^5} - \frac{a}{1+\dfrac{\tau^{(\mathrm{im})}}{\tau_0}\tilde{T}^3}\right)\sin\varphi\cos\varphi. \tag{26}$$

Here $a \equiv (\mu/\sigma_1)(\partial\sigma_1/\partial\mu) \cong 1$.

When $\sin 2\varphi = 0$, the magnitudes P_{xx}, P_{yy} do not depend on the magnetic field.

The nondiagonal components

$$P_{yx} = -\rho_0\tilde{T}\left[\sigma_2\left(\frac{1}{\omega_c\tau_p} - \frac{1-b}{\omega_c\tau_\varepsilon}\right) - \sigma_1\omega_c(\tau_p - a\tau_\varepsilon)\cos^2\varphi\right], \tag{27}$$

$$P_{xy} = \rho_0\tilde{T}\left[\sigma_2\left(\frac{1}{\omega_c\tau_p} - \frac{1-b}{\omega_c\tau_\varepsilon}\right) - \sigma_1\omega_c(\tau_p - a\tau_\varepsilon)\sin^2\varphi\right], \tag{28}$$

$$P_{zx} = \rho_0\tilde{T}\left[\sigma_2\left(\frac{d}{\omega_c\tau_p} - \frac{1-b}{\omega_c\tau_\varepsilon}\right) + \sigma_1\left(\omega_c(\tau_p - a\tau_\varepsilon)\cos^2\varphi + \frac{d}{\omega\tau_p}\right)\right]\tan\theta, \tag{29}$$

$$P_{zy} = \rho_0\tilde{T}\left[\sigma_2\left(-1 + (d+b)\frac{\tau_\varepsilon}{\tau_p}\right) + \sigma_1\left(\omega_c(\tau_p - a\tau_\varepsilon)\sin\varphi\cos\varphi + (d+b)\frac{\tau_\varepsilon}{\omega\tau_p}\right)\right]\tan\theta \tag{30}$$

also contain great terms proportional to γ^{-1}, the T-dependence of which is described by the expression analogous to formula (26).

It is easy to see that the Nernst–Ettingshausen effect is most pronounced when the temperature gradient is not orthogonal to the velocity $\mathbf{v}_1$.

The electric field along the normal to the layers grows with the angle θ increasing and exceeds the electric field along the layers at $\tan\theta > 1$.

When $\tan\theta > eB\tau/mc = 1/\gamma_0$, an electron has no time to perform a complete revolution along the closed cross-section of the Fermi surface during its free path time. An electron drifts to a small distance along the normal direction, and at $\theta = \pi/2$ its mean velocity along the normal to the layers equals to zero. As a result, in the main approximation in the quasi-two dimensionality parameter the conductivity component σ_{zz} is

$$\sigma_{zz} = \eta^2\gamma_0^2 g\sigma_2, \tag{31}$$

where g is the number of the order of unity. This is the case when the components of the tensor P_{ik}, determining the thermoelectric field along the normal to the layers,

$$P_{zx} = \tilde{T}\left[\omega_c\tau_p - \rho_0\sigma_2(1+b)\omega_c\tau_\varepsilon\right]\sin\varphi, \tag{32}$$

$$P_{zy} = -\tilde{T}\left[\omega_c\tau_p - (1+b)\omega_c\tau_\varepsilon\right]\cos\varphi, \tag{33}$$

grow linearly with the magnetic field increasing, while the components

$$P_{xx} = \rho_0\tilde{T}\left[\sigma_2(1+b) + \sigma_1(a+b)\right]\frac{\tau_\varepsilon}{\tau_p} \tag{34}$$

and

$$P_{yy} = \tilde{T}(1+b)\frac{\tau_\varepsilon}{\tau_p} \tag{35}$$

come up to saturation in the strong magnetic field. The electric field in the layers plane is directed mainly along the temperature gradient because the components

$$P_{yx} = P_{xy}(1+q) = \tilde{T}\eta^2 g \left[(1+b)\frac{\tau_\varepsilon}{\tau_p} - 1\right] \sin\varphi\cos\varphi \qquad (36)$$

are proportional to η^2 and, similarly to the diagonal components, tend to saturation at $\gamma_0 \ll 1$.

The presence of the charge carriers group with quasi-one-dimensional dispersion law does not influence markedly the thermoelectric effect when the magnetic field is almost parallel to the layers plane. If the vector **B** is deviated from the plane, both the T-dependence and the dependence on the magnetic field of the thermoemf prove to be essentially different for the conductor with one quasi-two-dimensional charge carriers group and for the conductor whose Fermi surface contains the plane sheets as well. The variety of these dependencies gives rich material for studying theproperties of charge carriers in low dimensional conductors and allows to reveal not only the presence of the plane sheets of the Fermi surface, but to determine the preferred direction of the velocity of electrons whose states belong to these sheets.

[1] O. V. Kirichenko, D. Krstovska, V. G. Peschansky, Sov. Phys. JETP **99**, 217 (2004).
[2] O. V. Kirichenko, V. G. Peschansky, R. A. Hasan, Low Temp. Phys. **32**, 1154 (2006).
[3] L. A. Falkovsky, Phys. Rev. B **75**, 033409 (2007).
[4] R. Rossenau *et al.*, J. Phys. I (France) **6**, 1527 (1996).
[5] H. Mori *et al.*, Bullet. Chem. Soc. Jpn **63**, 2183 (1990).

ЕЛЕКТРОННИЙ ТРАНСПОРТ У СИЛЬНО АНІЗОТРОПНИХ СТРУКТУРАХ У МАГНІТНОМУ ПОЛІ

О. В. Кириченко[1], В. Г. Піщанський[1], О. Галбова[2]
[1] *Фізико-технічний інститут низьких температур ім. Б. І. Веркіна Національної академії наук України, проспект Леніна, 47, Харків, 61103, Україна*
[2] *Факультет природничих наук і математики, фізичний іститут, Скоп'е, Македонія*

Теоретично досліджено лінійний відгук електронної системи шаруватого провідника на наявність ґрадієнта температури. Знайдено залежність термоедс від температури та зовнішнього магнітного поля, дослідження якої дасть змогу вивчити різні механізми релаксації в системі електронів провідності та визначити структуру електронного енерґетичного спектра.