Thermomagnetic phenomena in layered conductors

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A theoretical study is made of thermomagnetic phenomena in layered conductors in the presence of several groups of charge carriers. The thermopower at high external magnetic field is found as a function of the strength and orientation of the field; experimental study of this field dependence permits investigation of the structure of the energy spectrum of the charge carriers.
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Electronic phenomena in degenerate conductors at high magnetic fields **B**, when the frequency of gyration ω_c of the conduction electrons is much higher than their collision frequency $1/\tau$, are extremely sensitive to the form of the electron energy spectrum. Galvanomagnetic phenomena have been used successfully to recover the topology of the Fermi surface of metals from experimental study of the anisotropy of their magnetoresistance.[1,2] Thermomagnetic phenomena also contain rich information on the topological structure of the energy spectrum of the charge carriers.[3] In a quantizing magnetic field, when the width $\hbar\omega_c$ of the Landau levels is greater than the temperature smearing of the Fermi distribution function of the charge carriers, the thermoelectric field in low-dimensional conductors experiences giant oscillations as a function of the reciprocal of the magnetic field strength.[4] The condition of high magnetic field ($\omega_c\tau \gg 1$), necessary for solution of the inverse problem of recovering the electron energy spectrum from experimental data, is completely attainable in charge-transfer complexes having a layered structure.

Let us consider the thermomagnetic phenomena in layered conductors with an arbitrary sign of the dispersion of charge carriers at high magnetic field. Layered conductors are characterized by sharp anisotropy of their electrical conduction. The conductivity transverse to the layers is less than the conductivity along the layers by three orders of magnitude in organic conductors, by four orders of magnitude in manganites, and even by five orders of magnitude in graphite. This is apparently due to the sharp anisotropy of the charge-carrier velocities $\mathbf{v}=\partial\varepsilon(\mathbf{p})/\partial\mathbf{p}$, i.e., their energy $\varepsilon(\mathbf{p})$ depends weakly on the momentum projection $p_z=\mathbf{n}\cdot\mathbf{p}$ on the normal to the layers, **n**, and can be written in the form of a rapidly convergent series:

$$\varepsilon(\mathbf{p}) = \sum_{n=0}^{\infty} \varepsilon_n(p_x,p_y)\cos\left(\frac{anp_z}{\hbar}\right), \tag{1}$$

so that the projection of the electron velocity on the normal to the layers,

$$v_z = \frac{\partial\varepsilon}{\partial p_z} = -\sum_{n=1}^{\infty} \varepsilon_n(p_x,p_y)\frac{an}{\hbar}\sin\left(\frac{anp_z}{\hbar}\right) \tag{2}$$

is substantially lower than the maximum value of the velocity in the plane of the layers, v_F ($\hbar$ is Planck's constant, and a is the distance between layers). The maximum values of the functions ε_n on the Fermi surface fall off rapidly with increasing index n, and $\varepsilon_1^{\max}=\eta\varepsilon_F \ll \varepsilon_F$.

The Fermi surface $\varepsilon(\mathbf{p})=\varepsilon_F$ of a layered conductor is open and slightly corrugated along the p_z axis; it can be constructed with the aid of topologically simple elements in the form of slightly corrugated cylinders and corrugated planes, either isolated or connected pairwise by necks.

In many layered conductors the de Haas–van Alphen and Shubnikov–de Haas quantum oscillations have already been observed in different orientations of the magnetic field relative to the layers (see, e.g., the references cited in the review articles).[5,6] This attests to the fact that at least one sheet of the Fermi surface of these charge-transfer complexes is a slightly corrugated cylinder, since the closed electron orbitals on the corrugated planar sheets are utterly small, and almost the entire cross section $p_B=\mathbf{p}\cdot\mathbf{B}/B=$const of the planar sheets of the Fermi surface are open for any orientation of the magnetic field. Undoubtedly the Fermi surfaces of some layered conductors can consist of only a single slightly corrugated cylinder. In particular, it is customarily assumed that organic charge-transfer complexes based on tetrafulvalene, $(BEDT\text{–}TTF)_2X$ with $X=IBr_2$, I_3, have such a Fermi surface.[5] However, as a rule, the Fermi surface of layered conductors is single-sheet. There is reason to suppose that in the organic conductors $(BEDT\text{–}TTF)_2Cu(SCN)_2$ and $(BEDT\text{–}TTF)_2MHg(SCN)_4$, where M is one of the metals of the group K, Rb, Tl, or NH_4, there are two groups of charge carriers responsible for the charge transfer, which have a quasi-two-dimensional and a quasi-one-dimensional energy spectrum.[7] For interpretation of the phase relations between quantum oscillations of the magnetic susceptibility and magnetoresistance of graphite, Kopelevich and Luk'yanchuk[8] had to invoke three groups of carriers with topologically different character of the energy spectrum.

The linear response of the electron system to an external perturbation in the form of an electric field **E** and temperature gradient ∇T,

$$j_i = \sigma_{ij}E_j - \alpha_{ij}\frac{\partial T}{\partial x_j}, \tag{3}$$

$$q_i = \beta_{ij}E_j - \kappa_{ij}\frac{\partial T}{\partial x_j}, \tag{4}$$

can be found with the aid of a solution of the kinetic equation for the electron distribution function:

$$f(\mathbf{r},\mathbf{p}) = f_0(\varepsilon) - \psi_1(\mathbf{r},\mathbf{p})\frac{\partial f_0}{\partial \varepsilon} - \psi_2(\mathbf{r},\mathbf{p})\frac{\varepsilon-\mu}{T}\frac{\partial f_0}{\partial \varepsilon}, \quad (5)$$

where $f_0(\varepsilon)$ and μ are the equilibrium Fermi function and chemical potential of the electrons, T is the temperature in energy units, and the functions ψ_1 and ψ_2 are solutions of the equations

$$\frac{\partial \psi_1}{\partial t} + \mathbf{v}\frac{\partial \psi_1}{\partial \mathbf{r}} + \hat{W}_p\psi_1 = e\mathbf{E}\cdot\mathbf{v}, \quad (6)$$

$$\frac{\partial \psi_2}{\partial t} + \mathbf{v}\frac{\partial \psi_2}{\partial \mathbf{r}} + \hat{W}_\varepsilon\psi_2 = \mathbf{v}\frac{\varepsilon-\mu}{T}\nabla T. \quad (7)$$

Here e is the electron charge, the operators $\hat{W}_p$ and $\hat{W}_\varepsilon$ describe the relaxation of electrons with respect to the momenta $(1/\tau_p)$ and energies $(1/\tau_\varepsilon)$, and t is the time of motion of the charge in a magnetic field $\mathbf{B}=(B\cos\varphi\sin\theta, B\sin\varphi\sin\theta, B\cos\theta)$ according to the equations

$$\frac{\partial p_x}{\partial t} = \frac{eB\cos\theta}{c}(v_y - v_z\sin\varphi\tan\theta), \quad (8)$$

$$\frac{\partial p_y}{\partial t} = \frac{eB\cos\theta}{c}(v_z\cos\varphi\tan\theta - v_x), \quad (9)$$

$$\frac{\partial p_z}{\partial t} = \frac{eB\sin\theta}{c}(v_x\sin\varphi - v_y\cos\varphi). \quad (10)$$

The electric field and temperature gradient will be considered constant and uniform.

The eigenvalues of the scattering operators $\hat{W}_p$ and $\hat{W}_\varepsilon$ for conduction electrons, $1/\tau_p$ and $1/\tau_\varepsilon$, respectively, are substantially different, when the charge carriers are scattered on vibrations of the crystal lattice.[9,10] However, at low temperatures, once the condition $\omega_c\tau \gg 1$ is satisfied, the main mechanism of dissipation of the electrons at actually attainable magnetic fields is their scattering on impurity atoms and other crystal defects. The doping of a layered conductor by impurity atoms can substantially alter the electron energy spectrum.[11] We shall assume that the impurity centers are still too few in number to have a substantial influence on the energy spectrum of the charge carriers but entirely sufficient that we can neglect the electron-phonon scattering at low temperatures. In the scattering of charge carriers by impurity centers the momentum of the electron changes appreciably at the time of a collision event, and the relaxation times of the electron system with respect to momentum (τ_p) and energy (τ_ε) are of the same order of magnitude. In the case of a short range of interaction of the impurity center one can use the Born approximation for calculating the scattering amplitude. Here the collision integral, to a sufficient degree of accuracy, has the form of an operator of multiplication of the nonequilibrium correction to the distribution function by the collision frequency, and the solutions of equations (6) and (7) take the rather simple form

$$\psi_1 = \int_{-\infty}^{t} dt' \exp[(t-t')/\tau_p] e\mathbf{E}\cdot\mathbf{v}(t'), \quad (11)$$

$$\psi_2 = \int_{-\infty}^{t} dt' \exp[(t-t')/\tau_\varepsilon]\frac{\varepsilon-\mu}{T}\mathbf{v}(t')\nabla T. \quad (12)$$

After substituting the electron distribution function (5) into the expressions for the current density

$$\mathbf{j} = \frac{2}{(2\pi\hbar)^3}\int e\mathbf{v}f(\mathbf{r},\mathbf{p})d^3p \quad (13)$$

and heat flux density

$$\mathbf{q} = \frac{2}{(2\pi\hbar)^3}\int \mathbf{v}(\varepsilon-\mu)f(\mathbf{r},\mathbf{p})d^3p \quad (14)$$

one can easily determine the kinetic coefficients linking **j** and **q** with the electric field **E** and temperature gradient ∇T. From now on, we shall drop the distinction between τ_ε and τ_p and just set $\tau_\varepsilon = \tau_p = \tau$.

In the τ approximation for the collision integrals it is sufficient to calculate the components of the conductivity tensor σ_{ij}, while the remaining kinetic coefficients describing the heat transfer and thermoelectric effects are related to σ_{ij} by the simple relations

$$\alpha_{ij} = T^{-1}\beta_{ij} = \frac{\pi^2 T}{3e}\frac{\partial\sigma_{ij}(\mu)}{\partial\mu}, \quad (15)$$

$$\kappa_{ij} = \frac{\pi^2 T}{3e^2}\sigma_{ij}. \quad (16)$$

We shall assume that the Fermi surface consists of a corrugated cylinder and corrugated planes with an arbitrary corrugation along the p_y axis. The coordinate axes in the plane of the layers are directed so that the plane tangent to a corrugated planar sheet of the Fermi surface is parallel to the p_yp_z coordinate plane. In relations (13) and (14) it is necessary to integrate over all states of the conduction electrons, and in the presence of several groups of charge carriers, each of them brings its own contribution to the kinetic coefficients—in particular, to the components of the conductivity tensor, so that

$$\sigma_{ij} = \sigma_{ij}^{(1)} + \sigma_{ij}^{(2)}, \quad (17)$$

where $\sigma_{ij}^{(1)}$ is the contribution to the conductivity from charge carriers whose states are found on a planar sheet of the Fermi surface, and $\sigma_{ij}^{(2)}$ takes into account the contribution to σ_{ij} from the remaining electrons with the Fermi energy.

It follows from the equations of motion of a charge in a magnetic field that the charge carriers whose states belong to a corrugated planar sheet of the Fermi surface drift mainly along the x axis. After averaging Eq. (8) over a sufficiently long time segment, we find that the electron drift velocity along the y axis, $\bar{v}_y = \bar{v}_z\sin\varphi\tan\theta$ is of the same order as $\bar{v}_z$ if the magnetic field deviates substantially from the plane of the layers, i.e., if the angle Ψ is substantially different from $\pi/2$.

At the same time, a conduction electron can can move rather far along the p_y axis along its trajectory in **p** space, and, as follows from Eq. (9), the mean drift velocity $\bar{v}_x$ is of the order of v_F.

The presence of this group of charge carriers in addition to the conduction electrons on the corrugated cylinder of the

Fermi surface has a substantial influence on the dependence of the kinetic coefficients on the magnetic field strength at a high field. This is due to the fact that, however high the magnetic field, the high-field asymptote of the component $\sigma_{xx}^{(1)}(B)$ for $\gamma = T_B/\tau \ll 1$ is of the same order of magnitude as $\sigma_{xx}^{(1)}$ in the absence of magnetic field:

$$\sigma_{xx}^{(1)} = \frac{2e^2}{(2\pi\hbar)^3}\int d\varepsilon dp_y dp_z \frac{v_x^2\tau}{|v_x|}\delta(\varepsilon(\mathbf{p}) - \varepsilon_F)$$

$$= \frac{2e^2}{(2\pi\hbar)^3}\int dp_y dp_z |v_x|\tau = \frac{2e^2 v_1 \tau}{\pi\hbar ab}. \tag{18}$$

Here b is the period of the crystal lattice along the y axis, v_1 is the modulus of the mean drift velocity of the electrons along the x axis. The mean-free time for the two groups of electrons will be set equal, and the time T_B is of the same order of magnitude as the period of gyration of an electron along a closed orbit, while for electrons on an open trajectory in momentum space, it is of the same order of magnitude as the time of their displacement by a period of the reciprocal lattice.

In quasi-two-dimensional conductors the conductivity tensor components σ_{xx}, σ_{yy}, σ_{xy}, and σ_{yx} will be greater in order of magnitude than all the other components σ_{ij} even though they may decay with increasing magnetic field, since at actually attainable fields the small parameter γ is still much larger than the quasi-two-dimensionality parameter of the electron energy spectrum, i.e., $\eta \ll \gamma \ll 1$. These components are easily estimated by studying the Hall effect and magnetoresistance. In the leading approximation, the ratio of the electric fields E_y and E_x upon the flow of a current $\mathbf{j} = (j\cos\phi, j\sin\phi, 0)$ in the plane of the layers has the form

$$\frac{E_y}{E_x} = \frac{-\sigma_{xy}\cos\phi + \sigma_{xx}\sin\phi}{\sigma_{yy}\cos\phi - \sigma_{yx}\sin\phi}. \tag{19}$$

Having determined this ratio at four positions of current flow, e.g., at $\phi = 0$, $\pi/6$, $\pi/4$, and $\pi/2$, one can find the components of the conductivity tensor appearing in it for any form of the quasi-two-dimensional electron energy spectrum.

Further analysis of thermoelectric phenomena in conductors with a multi-sheet Fermi surface does not present difficulty. Upon a nonuniform heating of the conductor along the normal to the layers in the absence of current-conducting contacts ($\mathbf{j} = 0$), the thermoelectric field is directed mainly along the temperature gradient $\partial T/\partial z$:

$$E_z = \frac{\pi^2 T}{3e}\frac{\partial(\ln\sigma_{zz})}{\partial\mu}\frac{\partial T}{\partial z}. \tag{20}$$

The electric field in the plane of the layers, which is proportional to the square of the quasi-two-dimensionality parameter of the electron energy spectrum, though increasing with increasing magnetic field, is nevertheless, for $\eta \ll \gamma \ll 1$, much less than E_z:

$$E_x < E_y \simeq \eta^2 \frac{T}{\mu e}\tan\theta(\gamma^{-2}\sin\varphi + \gamma^{-1}\cos\varphi)\frac{\partial T}{\partial z}. \tag{21}$$

If the temperature gradient is oriented in the plane of the layers and directed along the y axis, then the thermoelectric field E_y is substantially larger than E_x and is given by the expression

$$E_y = \frac{\pi^2 T}{3e}\left(\rho_{yy}\frac{\partial\sigma_{yy}}{\partial\mu} + \rho_{yx}\frac{\partial\sigma_{xy}}{\partial\mu}\right)\frac{\partial T}{\partial y}$$

$$= \frac{\pi^2 T}{3e}\frac{\rho_0}{\sigma_{yy}^{(2)}}\left(-\sigma_{xx}^{(1)}\frac{\partial\sigma_{yy}^{(2)}}{\partial\mu} - \sigma_{yx}^{(2)}\frac{\partial\sigma_{xy}^{(2)}}{\partial\mu}\right)\frac{\partial T}{\partial y}, \tag{22}$$

where

$$\rho_0 = \frac{\sigma_{yy}^{(2)}}{\sigma_{yy}^{(2)}\sigma_{xx}^{(1)} - \sigma_{xy}^{(2)}\sigma_{yx}^{(2)}}.$$

The thermoelectric field

$$E_z = \frac{\pi^2 T}{3e}\frac{\rho_0}{\sigma_{yy}^{(2)}}\tan\theta\left\{(\sigma_{yx}^{(2)}\sin\varphi - \sigma_{yy}^{(2)}\cos\varphi)\frac{\partial\sigma_{xy}^{(2)}}{\partial\mu}\right.$$

$$\left. - (\sigma_{xx}^{(1)}\sin\varphi - \sigma_{xy}^{(2)}\cos\varphi)\frac{\partial\sigma_{yy}^{(2)}}{\partial\mu}\right\}\frac{\partial T}{\partial y} \tag{23}$$

is of the same order of magnitude as E_y only when the magnetic field deviates substantially from the normal to the layers, i.e., for $\tan\varphi \geqslant 1$.

The relation between the thermoelectric field and a temperature gradient directed along the x axis is given in the leading approximation in the parameter η as

$$E_x = \frac{\pi^2 T}{3e}\left(\rho_{xx}\frac{\partial\sigma_0^{(1)}}{\partial\mu} + \rho_{xy}\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right)\frac{\partial T}{\partial x}$$

$$= \frac{\pi^2 T}{3e}\rho_0\left(\frac{\partial\sigma_{xx}^{(1)}}{\partial\mu} - \frac{\sigma_{xy}^{(2)}}{\sigma_{yy}^{(2)}}\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right)\frac{\partial T}{\partial x}, \tag{24}$$

$$E_y = \frac{\pi^2 T}{3e}\left(\rho_{yx}\frac{3\sigma_0^{(1)}}{\partial\mu} + \rho_{yy}\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right)\frac{\partial T}{\partial x}$$

$$= \frac{\pi^2 T}{3e}\frac{\rho_0}{\sigma_{yy}^{(2)}}\left(-\sigma_{yx}^{(2)}\frac{\partial\sigma_{xx}^{(1)}}{\partial\mu} + \sigma_{xx}^{(1)}\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right)\frac{\partial T}{\partial x}, \tag{25}$$

$$E_z = \frac{\pi^2 T}{3e}\left\{\cos\varphi\frac{\partial}{\partial\mu}(\ln\sigma_{zz}) + \frac{\rho_0}{\sigma_{yy}^{(2)}}\left[\left(\sigma_{yx}^{(2)}\frac{\partial\sigma_{xx}^{(1)}}{\partial\mu}\right.\right.\right.$$

$$\left.\left.\left. - \sigma_{xx}^{(1)}\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right)\sin\varphi + \sigma_{xy}^{(2)}\cos\varphi\frac{\partial\sigma_{yx}^{(2)}}{\partial\mu}\right]\right\}. \tag{26}$$

The thermoelectric field turns out to be almost orthogonal to the temperature gradient. In the case of a Fermi surface consisting only of a single corrugated cylinder, the $\mathbf{E}$ vector lies mainly in the yz plane, and for $\theta \ll 1$ the component E_z is much less than E_x.

The contribution to the Hall component σ_{xy} from the charge carriers whose state belong to the calculated cylinder, viz.,

$$\sigma_{xy}^{(2)} = \frac{2ecS}{a(2\pi\hbar)^2 B} = \frac{N_2 ec}{B\cos\theta}, \tag{27}$$

is easily determined if one knows the period of the Shubnikov-de Haas oscillations, $\Delta(1/B) = 2\pi\hbar e/cS_{\text{extr}}$, since only this group of conduction electrons participates in the formation of the oscillations. In formula (27) S is the mean area of section of the Fermi surface by a plane p_B=const, which differs from the extremal cross sections of the Fermi surface, S_{extr}, by a negligible correction proportional to η.

Having determined σ_{xy} from experiment, we find

$$\sigma_{xy}^{(1)} = \sigma_{xy} - \sigma_{xy}^{(2)} = \frac{c}{eB}\int dp_B \int dp_y \frac{2}{(2\pi\hbar)^3}(\bar{p}_x - p_x(p_y)), \tag{28}$$

$\bar{p}_x$, as usual, is the value of the momentum component averaged over the time of motion of the charge. In the case of a slight corrugation of this sheet of the Fermi surface and along the p_y axis the Hall component $\sigma_{xy}^{(1)}$ will be negligible. Thus, knowing $\sigma_{xy}^{(1)}$, one can estimate the amount of corrugation of a planar sheet of the Fermi surface.

Apparently, $\sigma_{xy}^{(1)}$ is much less than $\sigma_{xy}^{(2)}$ even in the case when the corrugation of a planar sheet along the p_y axis is not small, and so

$$\frac{\partial \sigma_{xy}}{\partial \mu} \simeq \frac{\partial \sigma_{xy}^{(2)}}{\partial \mu} = \frac{4\pi ecm^*}{a(2\pi\hbar)^2 B}. \tag{29}$$

Thus, by studying the thermopower at high magnetic field, one can determine the carrier's cyclotron effective mass averaged over the Fermi surface, and from the value of $\sigma_{xx}^{(1)}$ determine the contribution to the conductivity of the sample from charge carriers whose states are found on a planar sheet of the Fermi surface.

For $\tan\theta \gg 1$ a substantial rearrangement of the electron trajectories in **p** space occurs. The sections through the corrugated cylinder by the plane p_B=const are so strongly elongated that an electron cannot complete a full orbit during the mean free time. In the case of a predominant direction of corrugation of a planar sheet of the Fermi surface, the dependence of the resistance on the strength of a magnetic field nearly orthogonal to the plane of a sheet is extremely peculiar.[2] If the corrugation of a planar sheet of the Fermi surface along the p_y axis is at least as large as the corrugation along the p_z axis, then its open cross sections are highly elongated along the p_z axis, when the magnetic field is almost orthogonal to the planar sheet of the Fermi surface. At $\theta=\pi/2$ the open electron trajectories in **p** space change direction, and the closed, highly elongated trajectories break up into pairs of trajectories open along the p_z axis, and the contributions to the conductivity from the two groups of charge carriers are of the same order of magnitude. Then the asymptotes of the tensor components $\sigma_{ij}^{(1)}$ at high magnetic field are of the same as the asymptotes of $\sigma_{ij}^{(2)}$, and the resistance to current transverse to the layers at magnetic fields that are not too high increases linearly with field for $\eta^{1/2} \ll \gamma \ll 1$, and with further increase of field for $\gamma \leqslant \eta^{1/2}$ the linear growth gives way to quadratic growth.[12] For the opposite relationship between the amount of corrugation of a planar sheet of the Fermi surface along the p_z and p_y axes, a substantial change of the magnetoresistance occurs as the magnetic field is rotated in the plane of the layers, when the **B** vector approaches the x axis. For $\varphi=0$ and $\theta=\pi/2$ the electrons whose states belong to a slightly corrugated cylinder drift in the plane of the layers, and along the z axis their motion is finite. Electrons whose states belong to a planar sheet of the Fermi surface, on the contrary, drift along the z axis. Thus, all the diagonal components of the total conductivity tensor for $\gamma \ll 1$ are nonzero, and the resistivity of the conductor for any orientation of the current density vector **j** is practically independent of the magnetic field strength and is of the same order of magnitude as in the absence of field.

The asymptotic expression for the tensor component $\sigma_{ij}^{(2)}$ at high magnetic field is not very sensitive to rotation of the magnetic field by a small angle in the plane of the layers, while at small φ the cross section of a planar sheet of the Fermi surface is highly elongated along the p_y axis, and for $\varphi \leqslant \gamma$ the electrons cannot travel a distance $2\pi\hbar/a$ along the p_z axis within the mean free time. The displacement of these electrons along the y axis is restricted in the same way as for $\varphi=0$. For $\varphi \gg \gamma$, however, the elongation of the electron trajectories along the p_y axis is no longer large, and the motion of the electrons turns out to be finite along the z axis. The asymptotes of the tensor components $\sigma_{ij}^{(1)}$ at high magnetic field in this region of angles ($\varphi \gg \gamma$) turns out to be the same as the asymptote of $\sigma_{ij}^{(2)}$, and the resistivity along the normal to the layer increases without bound with magnetic field. Thus the angle dependence of the magnetoresistance has a deep minimum at $\varphi=0$, the width of which, $\Delta\varphi \cong \gamma$, falls off with increasing magnetic field.

A study of the magnetoresistance of the organic conductors α-(BEDT–TTF)$_2$MHg(SCN)$_4$ and κ-(BEDT–TTF)$_2$Cu(SCN)$_2$ upon rotation of the magnetic field in the plane of the layers will permit an unambiguous determination of the extent to which one is justified in assuming that a group of charge carriers with a quasi-one-dimensional energy spectrum exists in them and to which the energy-band calculations done in Ref. 7 describe the real energy spectrum of the conduction electrons in these compounds.

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