Influence of optical alignment of momentum on galvanomagnetic effects in layer semiconductors

G. M. Shmelev, I. I. Zheru, Nguen Hong Shon, and G. I. Tsurkan

V. I. Lenin State University, Kishinev

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A theoretical investigation is made of the influence of optical alignment of the momentum of photocarriers on their kinetics in layer semiconductors subjected to nonquantizing magnetic fields. Quasielastic scattering is assumed in the case of nonequilibrium electrons transferred to the conduction band from the valence band by linearly polarized light. The β modification of a GaSe crystal is regarded as a typical layer semiconductor. The direct optical transition $\Gamma_4^- \to \Gamma_3^+$ in this semiconductor is forbidden (for a light beam perpendicular to the layer planes). Calculations are reported of the following "anomalous" transport effects: the transverse photoconductivity, longitudinal Hall effect, longitudinal magnetoresistance. An analysis is made of recent experiments which revealed a transverse photo-emf in GaSe.

Numerous investigations (see, for example, Refs. 1-12) of various aspects of the "anomalous" transport effects (transverse photo-emf, odd magnetoresistance, longitudinal magnetoresistance, transverse radioelectric effect, Hall effect in a longitudinal magnetic field, etc.) have established that they appear in initially optically isotropic semiconductors when these are subjected to polarized high frequency electric fields or electromagnetic waves. A necessary condition for the appearance of these effects is the anisotropy of the carrier momentum distribution induced by the incident polarized radiation. Specific mechanisms of manifestation of this anisotropy are various. They include, for example, the following: 1) a heating mechanism (associated with the intraband absorption of light)[1,2,4,5,9]; 2) a mechanism associated with the influence of a strong electromagnetic wave on the probability of the scattering of the band electrons by phonons or impurities[6,7]; 3) a photoionization mechanism in which the emitted electrons have a definite distribution of the emission angles and which is exhibited by extrinsic semiconductors[1,3]; 4) optical alignment (orientation) of photoelectron momenta in the course of interband absorption of light.[8,11,12]

We shall investigate the influence of the alignment of the momentum of photoelectrons on their kinetics in layer semiconductors subjected to nonquantizing magnetic fields. The scattering of nonequilibrium electrons transferred to the conduction band c from the valence band v by linearly polarized light is assumed to be quasielastic. The transport equation for the steady-state distribution function of electrons subjected to static electric and magnetic fields includes not only the field and collision terms, but also generation and recombination terms:

$$\left(e\mathbf{E} + \omega_H [\mathbf{p}, \mathbf{h}], \frac{\partial f(\mathbf{p})}{\partial \mathbf{p}}\right) = \mathrm{St}\,[f(\mathbf{p})] + G(\mathbf{p}) - Q(\mathbf{p}), \tag{1}$$

where $\mathrm{St}[f(\mathbf{p})]$ is the collision integral for electrons which we shall subsequently assume to be nondegenerate; $\mathbf{p}$ is the carrier momentum; $G(\mathbf{p})$ is the generation term; $Q(\mathbf{p})$ describes the loss of carriers from a band; $\omega_H = eH/mc$; m is the effective mass of a carrier in the conduction band; $\mathbf{h} = \mathbf{H}/H$.

We shall consider the direct optical transition $\Gamma_4^- \to \Gamma_3^+$ in layer semiconductors belonging to the crystal class D_{6h}. We shall assume that linearly polarized light is directed along the c axis of a crystal (we shall postulate that the polarization vector lies in the XOY plane: $\mathbf{F} \perp \mathbf{c}$). The matrix element for such transitions is $[\mathbf{e}, \mathbf{P}_{cv}(0)] = 0$ ($\mathbf{e} = \mathbf{F}/F$, where F is the amplitude of the electric field of the incident light). This follows directly from the symmetry properties of the Bloch functions $U_{c(v),\mathbf{p}}(\mathbf{r})$. In fact, if p = 0, we have

$$[\mathbf{e}, \mathbf{P}_{vc}(0)] \propto \left[\mathbf{e}, \int d\mathbf{r}\, U^*_{c,0}(\mathbf{r}) \nabla U_{v,0}(\mathbf{r})\right],$$

where the integrand transforms in accordance with the $\Gamma_3^+ \times \Gamma_5^- \times \Gamma_4^-$ representation that does not include a unit representation so that this integral vanishes [Γ_5 is a two-dimensional irreducible representation of the D_{6h} group governing the transformation of $\left(\frac{\partial}{\partial x} - i\frac{\partial}{\partial y}\right) - \left(\frac{\partial}{\partial x} + i\frac{\partial}{\partial y}\right)$]. Consequently, the $\Gamma_4^- \to \Gamma_3^+$ transition is forbidden. We then have

$$[\mathbf{e}, \mathbf{P}_{vc}(\mathbf{p})] = \mathrm{const}\,(\mathbf{p}, \mathbf{e}). \tag{2}$$

It therefore follows that in the case of the transition under consideration the generation term is given by $G(\mathbf{p}) \propto (\mathbf{p}, \mathbf{e})^2$. Expanding this expression in terms of spherical tensors (we are assuming that the Cartesian coordinate axes are oriented in an arbitrary manner), we obtain

$$G(\mathbf{p}) = G_0 (1 + b_{ik} \nu_i \nu_k)\, \delta(\varepsilon - \varepsilon_0), \tag{3}$$

where $b_{ik} = e_i e_k - (1/3)\delta_{ik}$; $G_0 = \mathcal{J} g^{-1}(\varepsilon_0) K_f$; $\mathcal{J}$ is the intensity of the incident light; $g(\varepsilon)$ is the density of states in the conduction band; K_f is the absorption coefficient for the forbidden transition.

The lifetime τ_c of an electron in the conduction band will be assumed to obey the isotropic relationship $\tau_c = \tau_c(\varepsilon)$, so that

$$Q(\mathbf{p}) = \frac{f(\mathbf{p})}{\tau_c(\varepsilon)}. \tag{4}$$

We shall seek a solution of Eq. (1) in the form of the expansion:

$$f(\mathbf{p}) = f^{(0)}(\varepsilon_p) + \frac{(\mathbf{f}^{(1)}, \mathbf{p})}{p} + \frac{f^{(2)}_{ik} p_i p_k}{p^2} \quad (i, k = 1, 2, 3). \tag{5}$$

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Following the procedure of Ref. 11, we obtain a system of equations that applies to any quasielastic carrier scattering mechanism:

$$\frac{1}{3p^2}\frac{\partial}{\partial p}p^2 e(\mathbf{E},\mathbf{f}^{(1)}) = I^{(0)}[f^{(0)}] + G_0\delta(\varepsilon-\varepsilon_0), \tag{6}$$

$$eE_i\frac{\partial f^{(0)}}{\partial p} + \omega_H\varepsilon_{ilk}h_l f_k^{(1)} + \frac{2}{5p^3}\frac{\partial}{\partial p}(p^3 e f_{ik}^{(2)}E_k) = I^{(1)}[f_i^{(1)}], \tag{7}$$

$$\left[p\frac{\partial}{\partial p}\left(\frac{eE_i f_k^{(1)}}{p} - \frac{1}{3}\frac{e(\mathbf{E},\mathbf{f}^{(1)})}{p}\delta_{ik}\right) + 2\omega_H\varepsilon_{ilm}h_l f_{mk}^{(2)}\right]_2 = I^{(2)}[f_{ik}^{(2)}] + G_2 b_{ik}\delta(\varepsilon-\varepsilon_0). \tag{8}$$

Here, $[\ldots]_2$ represents symmetrization of a tensor of rank 2, ε_{ikl} is a unit antisymmetric tensor,

$$I^{(l)}[f] = -\left(\frac{f}{\tau_c(\varepsilon)} + \frac{f}{\tau_{pl}(\varepsilon)}\right) + \frac{1}{g(\varepsilon)}\frac{\partial}{\partial\varepsilon} \times \left\{\frac{\varepsilon g(\varepsilon) f}{\tau_{\varepsilon l}(\varepsilon)} + \frac{\partial}{\partial\varepsilon}[D_{\varepsilon l}(\varepsilon)g(\varepsilon)f]\right\} \quad (l=0,1,2), \tag{9}$$

where $\tau_{pl}(\varepsilon)$ is the momentum relaxation time, $\tau_{\varepsilon l}(\varepsilon)$ is the energy relaxation time, and $D_{\varepsilon l}(\varepsilon)$ is the diffusion coefficient in the energy space (l labels the spherical harmonics of the distribution function). In the case of quasielastic scattering mechanisms we have $\tau_{pl} \ll \tau_{\varepsilon l}$, $T\varepsilon/D_{\varepsilon l}$ and, moreover, it is usually found that $\tau_{pl} \equiv \tau_{pl}(\varepsilon) \ll \tau_c$, so that we can assume

$$I^{(1,2)}[f^{(1,2)}] = -\frac{f^{(1,2)}}{\tau_{p1,2}}. \tag{10}$$

Equations (7) and (10) yield an expression for $\mathbf{f}^{(1)}(\varepsilon)$ in an approximation quadratic in E and then the current can be found from

$$\mathbf{j} = \frac{e}{m}\sum_{\mathbf{p}}\mathbf{p}f(\mathbf{p}) = \frac{e}{3m}\sum_{\mathbf{p}}p\mathbf{f}^{(1)}(\varepsilon_{\mathbf{p}}). \tag{11}$$

We thus obtain an expression for **j** which is valid for any quasielastic scattering mechanism:

$$\mathbf{j} = \frac{e^2 n}{m}\Big\{\langle\tau_{p1}\rangle\mathbf{E} - \omega_H\langle\tau_{p1}^2\rangle[\mathbf{h},\mathbf{E}] + \omega_H^2\langle\tau_{p1}^3\rangle(\mathbf{h}(\mathbf{h},\mathbf{E}) - \mathbf{E}) + \left\langle\!\left\langle\frac{\partial}{\partial\varepsilon}(\tau_{p1} - \omega_H^2\tau_{p1}^3)f^{(2)}\mathbf{E}\right\rangle\!\right\rangle - \omega_H\left[\mathbf{h},\left\langle\!\left\langle\frac{\partial}{\partial\varepsilon}\tau_{p1}^2 f^{(2)}\mathbf{E}\right\rangle\!\right\rangle\right] + \omega_H^2\mathbf{h}\left(\mathbf{h},\left\langle\!\left\langle\frac{\partial}{\partial\varepsilon}\tau_{p1}^3 f^{(2)}\mathbf{E}\right\rangle\!\right\rangle\right)\Big\}, \tag{12}$$

where

$$\langle\ldots\rangle = -\frac{2}{3n}\int_0^\infty(\ldots)g(\varepsilon)\frac{\partial f^{(0)}}{\partial\varepsilon}d\varepsilon, \quad \langle\!\langle\ldots\rangle\!\rangle = \frac{4}{15n}\int_0^\infty(\ldots)\varepsilon g(\varepsilon)\,d\varepsilon,$$

n is the carrier density, and $(f^{(2)}\mathbf{E})_i = f_{ik}^{(2)}E_k$. It should be noted that if τ_{p1} = const, then the photocurrent dependent on the polarization of the incident light [and represented by the last three terms in Eq. (12)] vanishes.

We shall now consider the case when $\tau_{p1,2}(\varepsilon) = \tau_{p1,2}(T)(\varepsilon/T)^{-r/2}$ and the equations for $f^{(0)}$ and $\hat{f}^{(2)}$ (in the approximation of zeroth order in E) are

$$0 = I^{(0)}[f^{(0)}] + G_0\delta(\varepsilon-\varepsilon_0); \tag{13}$$

$$[2\omega_H\varepsilon_{ilm}h_l f_{mk}^{(2)}]_2 = -\frac{f_{ik}^{(2)}}{\tau_{p2}} + G_0 b_{ik}\delta(\varepsilon-\varepsilon_0). \tag{14}$$

We shall find a solution of Eq. (13) on the assumption that only the "cold" carriers (of energy $\varepsilon \leq T$) are transferred from the conduction band and

$$\frac{1}{\tau_c(\varepsilon)} = \begin{cases} 0, & \varepsilon > T; \\ (T/\varepsilon)^{1/2}/\tau_c(T), & \varepsilon \leq T. \end{cases} \tag{15}$$

We shall also allow for the existence of dark carriers in the conduction band. We therefore obtain

$$f^{(0)}(x) = \frac{4\pi^{3/2}}{(2mT)^{3/2}}\exp(-x)\left[n_d + n_{ph} + \gamma n_{ph}\int_1^x\frac{\theta(x_0 - x')}{(x')^2}\exp(x')\,dx'\right], \tag{16}$$

where $x = \varepsilon/T$; n_d and n_{ph} are the densities of dark carriers and photocarriers, respectively; $\gamma = \tau_{ac}(T)/\tau_c(T)$; $\tau_{ac}(\varepsilon)$ is the energy relaxation time for the interaction with acoustic phonons. It should be pointed out that Eq. (16) satisfies the condition $f^{(0)}(\infty) = f^{(0)\prime}(\infty) = 0$ and also the requirement of equality of the number of photoelectrons transferred from the conduction band to the number of these electrons generated per unit time. For $x_0 \gg 1$ in the range $x \gg 1$ the function (16) is identical with the distribution function found in Ref. 13 specifically for this case.

We shall use Eq. (14) and iteration in respect of $\omega_H\tau_{p2}$ to find the function $f^{(2)}$ to within $(\omega_H\tau_{p2})^2$ and then, substituting this function together with Eq. (16) into Eq. (12), we shall obtain an expression for the current

$$j_i = \sigma_{ik}(H)E_k, \tag{17}$$

where

$$\sigma_{ik} = \sigma_0[\delta_{ik} - \lambda_1\omega_H\tau_{p1}(T)u_{ik} + \lambda_2\omega_H^2\tau_{p1}^2(T)\gamma_{ik}] - \sigma_2\,\mathrm{sign}\,r\,\{b_{ik} - \omega_H[(2\tau_{p10} + \tau_{p20})u_{im}b_{mk} - \tau_{p10}b_{im}u_{mk}] + \omega_H^2(3\tau_{p10}^2 + \tau_{p20}^2 + 2\tau_{p10}\tau_{p20})\gamma_{im}b_{mk} - 2(\tau_{p20}^2 + \tau_{p10}\tau_{p20})u_{im}b_{ms}u_{sk} + \tau_{p20}^2 b_{im}\gamma_{mk}]\}. \tag{18}$$

Here,

$$\lambda_1 = \frac{(n_d + n_{ph})\Gamma\left(\frac{5-2r}{2}\right) + \gamma_2 n_{ph}}{(n_d + n_{ph})\Gamma\left(\frac{5-r}{2}\right) + \gamma_1 n_{ph}}, \quad \lambda_2 = \frac{(n_d + n_{ph})\Gamma\left(\frac{5-3r}{2}\right) + \gamma_3 n_{ph}}{(n_d + n_{ph})\Gamma\left(\frac{5-r}{2}\right) + \gamma_1 n_{ph}}, \tag{19}$$

$$\sigma_0 = \frac{4}{3\sqrt{\pi}}\frac{e^2}{m}\tau_{p1}(T)\left[\Gamma\left(\frac{5-r}{2}\right)(n_d + n_{ph}) + \gamma_1 n_{ph}\right], \tag{20}$$

$$\sigma_2 = \sigma_0\frac{n_{ph}|r|x_0^{\frac{1-r}{2}}\tau_{p20}}{5\left[\Gamma\left(\frac{5-r}{2}\right)(n_d + n_{ph}) + \gamma_1 n_{ph}\right]\tau_\varepsilon(x_0)} \tag{21}$$

(σ_2 is the conductivity due to photoexcited carriers), $u_{ik} = \varepsilon_{ilk}h_l$, $\gamma_{ik} = h_i h_k - \delta_{ik}$,

$$\gamma_l = \gamma\left[\int_1^\infty x^{\frac{3-lr}{2}}e^{-x}dx\int_1^x\frac{\theta(x_0 - x')}{(x')^2}e^{x'}dx' - \int_1^{x_0}x^{-\frac{1+lr}{2}}dx\right] \equiv \gamma S_l(x_0); \tag{22}$$

for r = 1 (corresponding to the scattering by acoustic phonons) a calculation of the function $S_l(x_0)$ on a computer shows that if $1 < x_0 < 15$ then this function rises weakly and almost linearly from 0 to 0.93 (l = 1) and from 0 to 0.54 (l = 2) and then (in the range $x_0 > 15$) it remains practically constant: $S_1 \approx 0.97$ and $S_2 \approx 0.55$ ($S_3 \equiv 0$). If r =

-3 (scattering by ionized impurities), then

$$S_l \sim \frac{2}{3l-1} x_0^{\frac{3l-1}{2}} \quad \text{for} \quad x_0 > 10.$$

In the absence of a magnetic field, we find that

$$\mathbf{j} = \sigma_0 \mathbf{E} + \sigma_2 \operatorname{sign} r \left(\mathbf{e}(\mathbf{e}, \mathbf{E}) - \frac{1}{3} \mathbf{E} \right). \tag{23}$$

The photocurrent is given by

$$\mathbf{j}^{\text{photo}} = \sigma_2 \operatorname{sign} r \left(\mathbf{e}(\mathbf{e}, \mathbf{E}) - \frac{1}{3} \mathbf{E} \right). \tag{24}$$

It follows from Eq. (24) that, in particular, there is a transverse component of the current, first calculated for interband transitions in Ref. 8. The appearance of a transverse photo-emf in optically isotropic semiconductors was predicted in Refs. 1 and 2 on the basis of general (group-theoretic) considerations. According to Refs. 1 and 2, the photocurrent is

$$j_{\parallel}^{\text{photo}} = \left[\frac{1}{2}(\gamma_{11} + \gamma_{12}) + \frac{1}{2}(\gamma_{11} - \gamma_{12}) \cos 2\varphi \right] EF^2,$$

$$j_{\perp}^{\text{photo}} = \gamma_{44} \sin 2\varphi EF^2.$$

In the case of the specific mechanism of a transverse emf considered here, we have $\gamma_{11} = 2\sigma_2 \operatorname{sign} r/3F^2$; $\gamma_{12} = -\sigma_2 \operatorname{sign} r/3F^2$, $\gamma_{44} = \sigma_2 \operatorname{sign} r/2F^2$ and as expected, $\gamma_{44} = \frac{1}{2} \cdot (\gamma_{11} - \gamma_{12})$.

We shall assume that linearly polarized light travels along the OZ axis and that a longitudinal field is directed along the OX axis (we shall assume that the sample is open-circuited along the OX and OZ directions).

A. H || OZ. The transverse magnetoresistance is

$$\left(\frac{\Delta\rho}{\rho}\right)_{\perp} = -\frac{\sigma_2 \operatorname{sign} r}{\sigma_0} \omega_H (\tau_{p10} + \tau_{p20}) \sin 2\varphi + \omega_H^2 \Big\{ \tau_{p1}^2(T) \Big[\lambda_2 - \lambda_1^2 - \frac{\sigma_2 \operatorname{sign} r}{3\sigma_0} \times (\lambda_2 (2 - 3\sin^2\varphi) - \lambda_1^2) \Big] - \frac{\sigma_2 \operatorname{sign} r}{\sigma_0} \Big[\frac{2}{3} \lambda_1 \tau_{p1}(T) \tau_{p10} - 3\tau_{p10}^2 \cos^2\varphi - 2(\tau_{p20}^2 + \tau_{p10}\tau_{p20}) \cos 2\varphi \Big] \Big\}, \tag{25}$$

where φ is the angle between **F** and the OX axis. The static transverse field is

$$E_y = -\frac{j\sigma_2 \operatorname{sign} r}{2\sigma_0^2} \sin 2\varphi + jHR_{\perp}(H). \tag{26}$$

The Hall coefficient is given by

$$R_{\perp}(H) = R_{\perp} \Big\{ 1 - \frac{\sigma_2 \operatorname{sign} r}{\sigma_0} \Big[\frac{1}{3} - \omega_H \tau_{p1}(T)(\lambda_1 - 2\lambda_2/\lambda_1) \frac{\sin 2\varphi}{2} - \frac{\omega_H \sin 2\varphi}{2\tau_{p1}(T)\lambda_1} (3\tau_{p10}^2 + 4\tau_{p20}^2 + 4\tau_{p10}\tau_{p20}) - \frac{2\tau_{p10}(\cos^2\varphi - 1/3) + \tau_{p20}\cos 2\varphi}{\tau_{p1}(T)\lambda_1} \Big] \Big\}, \tag{27}$$

where

$$R_{\perp} = \frac{3\sqrt{\pi}}{4ec} \frac{\lambda_1}{(n_d + n_{ph})\Gamma\left(\frac{5-r}{2}\right) + \gamma_1 n_{ph}}. \tag{28}$$

B. H || OX ($x_0 \gg 1$). The current along the OX axis is

$$j_x = \sigma_0 E \Big\{ 1 + \frac{\sigma_2 \operatorname{sign} r}{\sigma_0} (\cos^2\varphi - 1/3) - \left(\frac{\sigma_2 \sin 2\varphi}{2\sigma_0} \right)^2 \times [1 + \omega_H^2 (\tau_{p1}^2(T)(\lambda_2 - \lambda_1^2) + \tau_{p1}(T)\lambda_1(\tau_{p10} + \tau_{p20}) - 3\tau_{p10}^2)] \Big\}. \tag{29}$$

The transverse field is

$$E_y = -\frac{j\sigma_2 \operatorname{sign} r \sin 2\varphi}{2\sigma_0^2} \{ 1 + \omega_H^2 [\tau_{p1}^2(T)(\lambda_2 - \lambda_1^2) + \tau_{p1}(T)(2\tau_{p10} + \tau_{p20})\lambda_1 - (3\tau_{p10}^2 + \tau_{p20}^2 + 2\tau_{p10}\tau_{p20})] \}, \tag{30}$$

$$E_z = jHR_{\parallel},$$

where

$$R_{\parallel} = -R_{\perp} \frac{\sigma_2 \operatorname{sign} r}{2\sigma_0} \sin 2\varphi \left[1 - \frac{2\tau_{p10} + \tau_{p20}}{\lambda_1 \tau_{p1}(T)} \right]. \tag{31}$$

Before considering the characteristic features of the dependences (25)-(31), we must point out that the results obtained above apply not only to the case of interband excitation of carriers, but also in general to any excitation mechanism which can be described by Eq. (3).

Linearly polarized light (and a nonequilibrium of the electron system) stimulate an odd (in respect of the magnetic field) magnetoresistance represented by the first term in Eq. (25). This odd magnetoresistance is entirely due to nonequilibrium carriers during the ballistic stage of their motion in the conduction band, i.e., from the moment of their creation to the moment of loss of directional momentum. The odd magnetoresistance is characterized also by a special angular dependence $\sin 2\varphi$ which can be used to identify this effect experimentally. The sign of the odd magnetoresistance depends also on the sign of the quantity representing the nature of quasielastic scattering. It should be pointed out that attention to the possibility of the appearance of an odd interband photomagnetoresistance effect was first drawn in Ref. 11, where the scattering by acoustic phonons in centrosymmetric semiconductors was considered.

It follows from Eq. (26) that, in particular, a transverse photo-emf differs from zero even in H = 0. Recently a transverse photo-emf was observed experimentally in GaSe crystals (of the β modification) as a result of the $\Gamma_4^- - \Gamma_3^+$ interband transition.[14] An analysis of the results reported in Ref. 14 by means of Eq. (26) demonstrates an anomalously high mobility u_{ph} of photoexcited optically aligned carriers compared with the mobility of thermalized and dark carriers (u_0): $u_{ph} \sim 10^2 u_0$. This shows that optical alignment of the carrier momenta and ballistic motion are manifested in these experiments.

The expressions given in Eq. (30) are responsible for another "anomalous" effect: the Hall effect due to photocarriers in a longitudinal magnetic field. It was first considered in Ref. 12 in the specific case of the scattering by acoustic phonons and forbidden transitions. The expression (29) for the current j_x allows for the possibility of existence of a photostimulated interband longitudinal magnetoresistance:

$$\left(\frac{\Delta\rho}{\rho}\right)_{\parallel} = \left(\frac{\Delta\rho}{\rho}\right)_{\perp} \left(\frac{\sigma_2}{2\sigma_0} \sin\varphi \right)^2 \frac{r^2}{\lambda_1^2}, \tag{32}$$

where $(\Delta\rho/\rho)_{\perp}$ is the "usual" even transverse magnetoresistance. The dependence (32) can be explained as follows. A polarized electromagnetic wave stimulates the appearance of a static electric field along the OY

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axis. The presence of a magnetic field along the OX axis "rotates" this field (E_y) in the direction of the OZ axis. The return "rotation" in a magnetic field toward the OY axis and in the field of the electromagnetic wave to the OX axis alters the initial current, i.e., it gives rise to the magnetoresistance. Since each "rotation" in a magnetic field corresponds to a factor $\sim(\Delta\rho/\rho)_H^{1/2}$, and that in the field of the electromagnetic wave corresponds to $\sim\sigma_2/2\delta_0 \sin 2\varphi$, these considerations yield (apart from a numerical factor) Eq. (32), found by direct calculation.

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Two-particle tunneling in a normal metal–semiconductor contact

A. L. Shelankov

A. F. Ioffe Physicotechnical Institute, Academy of Sciences of the USSR, Leningrad

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It is shown that at low temperatures the resistance of a tunnel contact between a normal metal and a semiconductor is determined by two-particle tunneling (Andreev reflection). The ohmic resistance at $T \ll T_c$ and the excess current are calculated. A theory of Andreev reflection at a boundary with arbitrary transmittance is constructed.

The current–voltage characteristic of a tunnel NIS contact formed by a normal metal (N) and a superconducting metal (S), separated by a layer of dielectric (I), was first calculated by Cohen et al.[1] Using the tunneling Hamiltonian method the contribution to the current through the contact made by tunneling of single-particle excitations was obtained in Ref. 1. At low voltages in single-particle tunneling only excitations with energy greater than the superconductor gap participate. At low temperatures $T \ll \Delta$ the number of such excitations is exponentially small, and the ohmic conductance of the contact, calculated in the single-particle tunneling approximation, vanishes as $T \to 0$ as $\exp(-\Delta/T)$.[1] This indicates the low-efficiency of single-particle processes in tunnel conductance at low temperatures.

One other mechanism exists for charge exchange between N and S metals: the transfer of two electrons of the N-metal with energies $\mu_S + \xi$ and $\mu_S - \xi$ into the S metal with the formation of a Cooper pair (μ_S is the electrochemical potential of S), and also the inverse transitions. For the NS boundary this process is known as Andreev reflection,[2] whereas in the case of an NIS contact it must be regarded as two-particle tunneling. The efficiency of two-particle processes is determined by the following considerations. The probability of two electrons tunneling, which is proportional to the square of the tunnel transmittance, is small compared with the probability of single-particle tunneling. On the other hand, the excitation energy ξ of the electrons of the N-metal which participate in two-particle tunneling may be as small as desired, i.e., this process does not require activation. We would therefore expect that at fairly low temperatures of nonactivation two-particle tunneling processes, despite their low probability, would be a more effective mechanism of tunnel conductance than single-particle processes.

The purpose of the present paper is to set up a theory of the NIS contact taking two-particle process into account.[1)]

The single-particle tunnel current can be expressed solely in terms of the density of states, and does not depend explicitly on such properties of the metals as the mean free path. It is not possible to construct such a universal theory of two-particle processes. In this paper we only consider the case of a long mean free path – a clear limit. In the clear limit, as will be shown below, to calculate the current through the contact it is sufficient to obtain the probability of different excitation scattering channels at the NIS boundary, taking into account the finite transmittance of the tunnel layer s.

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