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Abstract

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PHYSICS

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ON THE QUANTUM THEORY OF THE PHOTOMAGNETIC EFFECT IN SEMICONDUCTORS

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1. The experimental investigations of photomagnetic effects in semiconductors that have been carried out intensively in recent years, covering the region $\Omega\tau \gg 1$, nevertheless pertained to the interval of temperatures and fields for which $\hbar\Omega \ll kT$ ($\Omega = eH/mc$ is the Larmor frequency of a quasiparticle, τ is the relaxation time in the absence of a magnetic field). This made it possible, in a theoretical study of photomagnetic effects in semiconductors, to restrict oneself to a quasiclassical treatment based on the usual kinetic equation (see, for example, ([1])). However, for photomagnetic effects the study of the quantum region ($\hbar\Omega \gg kT$) is of special interest, since the effects themselves depend to a decisive degree on the simultaneous motion of two types of carriers. The presence of electrons and holes with substantially different effective masses, and even obeying different statistics, predetermines the occurrence of a number of cases in which a quantum treatment will lead to completely different types of dependences of the electric field E_1, arising as a result of the photomagnetic effect, on the magnetic field and on the parameters of the semiconductor. To find this field in the case under consideration it is necessary to use the known quantum-mechanical description of particle motion in a strong magnetic field and, at the same time, the density-matrix apparatus. This problem is solved in the present work for the case of an unsteady photomagnetic effect under the assumption of isotropy of the dispersion laws for quasiparticles of both types and of an impurity scattering mechanism.

2. Let us consider a semiconductor in the form of a plane-parallel plate of thickness d along axis 2, in the direction of which the carriers generated by light diffuse. Let the magnetic field be directed along axis 3. Then the electric field of the photomagnetic effect, directed along axis 1, will be determined from the solution of the system of equations (see, for example, ([1]))

$$\int_0^d (\mathbf{n}_1, \mathbf{j})\, dx_2 = 0, \qquad (\mathbf{n}_2, \mathbf{j}) = 0, \qquad \operatorname{rot} \mathbf{E} = 0, \tag{1}$$

where $\mathbf{j}$ is the total current density, $\mathbf{n}_i$ are the unit vectors of the coordinate system. As direct estimates show (see, for example, ([2])), the thermalization time of the carriers even at very low temperatures is small compared with the lifetime. This makes it possible to calculate the current as arising only at constant temperature, equal to the lattice temperature, from the density gradient and the electric field. In doing so we may make use of the results obtained for the diamagnetic and dissipative currents of particles in a strong magnetic field in the presence of spatial inhomogeneity ([3,4])

$$j_\alpha = j_{\alpha 0} + j_{\alpha 1},$$

$$j_{\alpha 0} = \mathbf{n}_1 \left(\frac{\partial \mu_\alpha}{\partial x_2} \right) \left\{ -\beta |e_\alpha| \Omega_\alpha \sum_\nu f_\nu^\alpha (1 - f_\nu^\alpha) \left(n + \frac{1}{2} \right) + \frac{n_\alpha^0 c}{H} \right\}, \tag{2}$$

$$j_{\alpha 1}^i = \sigma_\alpha^{ik} \left(E_k - \frac{1}{e_\alpha} \frac{\partial \mu_\alpha}{\partial x_k} \right), \qquad i, k = 1, 2.$$

Here f_ν^α is the Fermi or Boltzmann distribution function in the state $\nu = (p_1, p_3, n)$,

$$\begin{aligned} \sigma_\alpha^{12} &= -\sigma_\alpha^{21} = \frac{c e_\alpha n_\alpha^0}{H}, \qquad n_\alpha^0 = \sum_\nu f_\nu^\alpha, \\ \sigma_\alpha^{11} &= \sigma_\alpha^{22} = \frac{e_\alpha^2 \beta}{2} \int_{-\infty}^{\infty} dt \, \langle \hat{X}_\alpha(0) \hat{X}_\alpha(t) \rangle, \end{aligned} \tag{3}$$

where $\beta = 1/kT$, and $\hat{X}_\alpha$ is the operator of the center of the Larmor motion of the quasiparticle. For σ_α^{11} one may use the results obtained by taking into account scattering by uncorrelated impurities, using the exact scattering amplitude on a short-range potential in a strong magnetic field ([5,6]).

When substituting (2) into (1), it is also necessary to take into account the existence of a surface current arising in a spatially inhomogeneous system placed in a quantizing magnetic field ([7,8]), and strictly compensating the diamagnetic current $j_{\alpha 0}$. As a result, considering for definiteness an n-type semiconductor (for which the inequality $n_0 \gg p_0$ holds, where n_0, p_0 are the equilibrium concentrations of electrons and holes, respectively), we find

$$E_1 = -\frac{1}{e_p} \frac{[\mu_p(0) - \mu_p(d)]}{d} \frac{\sigma_p^{12} \sigma_n^{22} - \sigma_p^{22} \sigma_n^{12}}{\sigma_n^{12} \sigma_n^{21} - \sigma_n^{11} \sigma_n^{22}}. \tag{4}$$

Using (3) and the inequality $\left|\sigma_n^{12}/\sigma_n^{11}\right| \gg 1$, and also assuming that, because of their small concentration, the holes are always described by the Boltzmann distribution $f_\nu^p = \exp[-\beta(\varepsilon_\nu^p - \mu_p)]$, from (4) we obtain

$$E_1 = A_p \frac{H}{ec}\left(\frac{\sigma_n^{22}}{n_0} + \frac{\sigma_p^{22}}{p_0}\right), \tag{5}$$

where

$$A_p = -\frac{1}{en_0\beta}\frac{\Delta p(0) - \Delta p(d)}{d}, \qquad e = |e_\alpha|, \qquad \Delta p \ll p_0.$$

3. Let us analyze formula (5) in various regions with respect to temperature and magnetic field.

I. We begin with the case when the electronic component of the current carriers is degenerate, i.e.,

$$f_\nu^n = \{\exp[\beta(\varepsilon_\nu^n - \mu_n^0)] + 1\}^{-1}.$$

1) Let the magnetic field be so strong that

$$\hbar\Omega_p\beta \gg 1, \qquad \hbar\Omega_n\beta \gg 1, \qquad \hbar\Omega_n/\mu_n^0 \gg 1.$$

Then from (5) we find:

$$E_1 = \frac{A_p}{4}\left\{\frac{27\pi}{32}\left(\frac{\hbar\Omega_n}{\mu_n^0}\right)^5 \frac{1}{\Omega_n\tau_n^F} + \frac{(\hbar\Omega_p\beta)^2}{\Omega_p\tau_p^0}\left[-\operatorname{Ei}\left(-\frac{f_p^2}{\lambda_H^2}\frac{\hbar\Omega_p\beta}{2}\right)\right]\right\}. \tag{6}$$

Here f_p is the exact scattering amplitude of a hole, $\lambda_H^2 = c\hbar/eH$, $(\tau_\alpha^0)^{-1} = N_{\text{imp}}(4\pi f_\alpha^2)\bar{v}_\alpha$, $(\tau_\alpha^F)^{-1} = N_{\text{imp}}(4\pi f_\alpha^2)v_\alpha^F$, where N_{imp} is the number density of impurities, and $\bar{v}_\alpha$ and v_α^F are, respectively, the mean thermal velocity and the velocity at the Fermi surface.

It follows from (6) that when the first term is small, everything is determined by the holes, and the dependence on the magnetic field has the form

$$E_1 \sim H \ln H^2.$$

When, however, the main role is played by electrons, then

$$E_1 \sim H^4.$$

However, obtaining such strong magnetic fields experimentally is very difficult, so that the following case is more realistic.

2) Consider the case when

$$\hbar\Omega_p\beta \gg 1, \qquad \hbar\Omega_n\beta \gg 1, \qquad \hbar\Omega_n/\mu_n^0 \ll 1.$$

In this case

$$E_1 = \frac{A_p}{2}\left\{\frac{1}{\Omega_n\tau_n^F} + \frac{(\hbar\Omega_p\beta)^2}{2(\Omega_p\tau_p^0)}\left[-\operatorname{Ei}\left(-\frac{f_p^2}{\lambda_H^2}\frac{\hbar\Omega_p\beta}{2}\right)\right]\right\} + E_{\text{osc}}, \tag{7}$$

where

$$E_{\text{osc}} = \frac{5}{4\sqrt{2}}\frac{A_p}{(\Omega_n\tau_n^F)}\sqrt{\frac{\hbar\Omega_n}{\mu_n^0}}\exp\left(-\frac{2\pi^2}{\hbar\Omega_n\beta}\right)\cos\left(2\pi\frac{\mu_n^0}{\hbar\Omega_n} - \frac{\pi}{4}\right), \tag{8}$$

and, naturally, in (8) we have retained only the principal term of the sum. Thus, on a monotonic dependence of the form

$$E_1 \sim H^{-1} \quad \text{or} \quad E_1 \sim H\ln H^2$$

typical quantum oscillations will be observed as H passes through values determined by the relation $\hbar\Omega_n M = \mu_n^0$, where M is an integer.

3) Further, the case is possible

$$\hbar\Omega_p\beta \ll 1, \qquad \hbar\Omega_n\beta \gg 1, \qquad \hbar\Omega_n/\mu_n^0 \ll 1,$$

which, evidently, corresponds to the inequality $m_p \gg m_n$. Then

$$E_1 = A_p\left[\frac{1}{\Omega_p\tau_p^0} + \frac{1}{2(\Omega_n\tau_n^F)}\right] + E_{\text{osc}}, \tag{9}$$

where the monotonic dependence $E_1 \sim H^{-1}$ is determined here, in contrast to (7), by both electrons and holes.

4) Finally, if

$$\hbar\Omega_p\beta \ll 1, \qquad \hbar\Omega_n\beta \ll 1,$$

then we arrive at the quasiclassical limit

$$E_1 = A_p \left[\frac{1}{\Omega_p \tau_p^0} + \frac{1}{2(\Omega_n \tau_n^F)} \right]. \tag{10}$$

II. Let us now consider the case when the electrons, like the holes, are described by the Boltzmann distribution

$$f_\nu^n = \exp\left[-\beta(\varepsilon_\nu^n - \mu_n^0)\right].$$

In this case, different relations between the temperature and the magnetic field are again possible.

1) If

$$\hbar\Omega_p\beta \gg 1, \qquad \hbar\Omega_n\beta \gg 1,$$

then

$$E_1 = -\frac{A_p}{4}\left\{ \frac{(\hbar\Omega_p\beta)^2}{\Omega_p\tau_p^0}\,\mathrm{Ei}\left(-\frac{f_p^2}{\lambda_H^2}\frac{\hbar\Omega_p\beta}{2}\right) + \frac{(\hbar\Omega_n\beta)^2}{\Omega_n\tau_n^0}\,\mathrm{Ei}\left(-\frac{f_n^2}{\lambda_H^2}\frac{\hbar\Omega_n\beta}{2}\right)\right\}. \tag{11}$$

In this case, for any relation between the first and second terms,

$$E_1 \sim H \ln H^2.$$

2) When, however,

$$\hbar\Omega_p\beta \ll 1, \qquad \hbar\Omega_n\beta \gg 1,$$

then

$$E_1 = A_p/\Omega_p\tau_p^0. \tag{12}$$

Thus, in this case the quasiclassical limit is determined by the carriers with the heavier effective mass.

3) If

$$\hbar\Omega_p\beta \ll 1, \qquad \hbar\Omega_n\beta \ll 1,$$

then we have a result coinciding with the limiting value of the emf of the photomagnetic effect obtained by the method using the kinetic equation:

$$E_1 = A_p \left[\frac{1}{\Omega_p \tau_p^0} + \frac{1}{\Omega_n \tau_n^0} \right], \tag{13}$$

when

$$E_1 \sim H^{-1}.$$

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Note: Figure translations are in progress. See original paper for figures.