Alternating oscillations in the photocurrent spectra of nonequilibrium photoelectrons in p-type InSb subjected to quantizing magnetic fields

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(Submitted July 11, 1972)
Zh. Eksp. Teor. Fiz. **64**, 231-244 (January 1973)

An investigation was made of the photocurrent spectra of Ge-doped p-type InSb subjected to quantizing magnetic fields at 4.2°K. A reversal of the sign of the transverse photocurrent was observed in magnetic fields in excess of 18 kOe when the photoelectrons were excited to the lower Landau levels. Moreover, photocurrent minima were observed due to the de-excitation of photoelectrons to the lower Landau levels because of the emission of optical phonons. These minima were located in the positive region; they were shallow and much wider than the minima associated with the excitation of photoelectrons to the Landau levels. A study was made of the dependence of the negative photocurrent on the monochromaticity and the intensity of the incident radiation. The experimental observations of the negative photocurrent and the good agreement between the experimentally and theoretically obtained spectra indicated that the absolute negative conductivity of photoelectrons was observed when the investigated samples of p-type InSb were subjected to quantizing magnetic fields at helium temperatures.

1. INTRODUCTION

Several mechanisms of negative photoconductivity in semiconductors have been investigated experimentally (see, for example,[1,2]). Usually a negative photoconductivity appears as a result of a reduction in the carrier mobility or density but not as a result of a change in the direction of motion of electrons. The negative photoconductivity due to an anomalous direction of the motion of charges relative to the electric field is of special interest. In the presence of this absolute negative conductivity one should be able to observe, at least in principle, the flow of the total current against the electric field as well as various instabilities.

In earlier papers[3] Elesin predicted an absolute negative conductivity of nonequilibrium electrons in quantizing magnetic fields due to the special quantization-induced dependence of the density of states on the energy. In the presence of scattering an electron subjected to mutually perpendicular magnetic and electric fields can travel along the electric field only if its momentum along the magnetic field is greater than a certain critical value (apart from that the electron still drifts along the third direction, in the usual manner). This occurs because in the elastic or almost elastic scattering of an electron with an initial momentum $|p_z| > p_{cr} \propto (2me\mathscr{E}\Lambda)^{1/2}$ the density of the final states for jumps along the electric field, proportional to $[(p_z^2/2m) - e\mathscr{E}\Delta X]^{-1/2}$, exceeds the density of the final states for jumps against the field, proportional to $[(p_z^2/2m) + e\mathscr{E}\Delta X]^{-1/2}$, and the jumps along the field are most probable. Here, $\mathscr{E}$ is the electric field, e is the absolute value of the charge, $\Delta X \propto \Lambda = (eH)^{-1/2}$ is the length of a jump (for the sake of simplicity we shall consider an electron at a lower Landau level and we shall assume that $\hbar = c = 1$).

Electrons located in the direct vicinity of the lower edges of the Landau sub-bands with $|p_z| < p_{cr}$ can only jump against the field because the density of the final states in jumps of these electrons along the field is zero. In a weak electric field the quantity p_{cr} can be shown to be proportional to $(2m\Gamma)^{1/2}$, where Γ is the collision broadening of the Landau levels. In the case of an equilibrium distribution the electrons with $|p_z| < p_{cr}$ always make a contribution which is larger than the negative contribution of the electrons with $|p_z| > p_{cr}$ and the current flows along the field. An absolute negative conductivity appears only in the presence of a distribution with a positive derivative with respect to the energy near the lower edges of the Landau sub-bands. This distribution can be established if the photoelectron lifetime in an allowed band is less than the energy relaxation lifetime, $\tau_R < \tau_\epsilon$, and if the photoelectrons are created with momenta $|p_z| > p_{cr}$. The possibility of establishing a strongly nonequilibrium distribution of photoelectrons in p-type InSb was demonstrated experimentally by Habegger and Fan.[4]

The present authors[5] have suggested that measurements of the spectral distribution of the photocurrent in quantizing magnetic fields provides a method for the observation of the absolute negative conductivity. Oscillations of the photocurrent observed for some p-type InSb samples have been interpreted successfully on the basis of the absolute negative conductivity of electrons far from equilibrium. However, a negative photocurrent is not observed because of the presence of a strong positive background of thermalized photoholes (this should be compared with the oscillations of the photocurrent in GaSb reported in[6]).

The present paper reports a study of the negative photoconductivity of p-type InSb observed in quantizing magnetic fields at liquid helium temperature because of the absolute negative photoconductivity of electrons far from equilibrium. Detailed experimental and theoretical studies of the spectral oscillations of the photocurrent are also reported.

2. EXPERIMENTAL RESULTS

It follows from the Introduction that the necessary condition for the appearance of an absolute negative conductivity is a nonequilibrium energy distribution of photoelectrons. Such nonequilibrium conditions can be established in germanium-doped p-type InSb. According to Berkeliev et al.,[7] the minority carrier lifetime in

this material is sufficiently short even at liquid nitrogen temperature ($\tau_R \lesssim 10^{-9}$ sec) and this lifetime has a tendency to decrease with increasing temperature.

We determined the photocurrent spectra of a large batch of samples (over 40) in which the majority-carrier density was $p_0 \approx 10^{13}$ cm^{-3} and the mobility was $\mu_p = (4-8) \times 10^3$ cm$^2 \cdot$ V$^{-1} \cdot$ sec^{-1} (at 77° K). Measurements of the Hall effect indicated that the majority-carrier mobility at 4.2° K was low: $\mu_p \approx 50-250$ cm$^2 \cdot$ V$^{-1} \cdot$ sec^{-1}. Our samples were 7—15 mm long, 1—3 mm wide, and 0.3—2 mm thick. The thickness of the samples exceeded slightly the diffusion and the absorption lengths.

After careful polishing and etching in CP-4A a sample was placed in a superconducting solenoid of the Helmholtz-coil type in which magnetic fields up to 50 kOe could be generated. The measurements were carried out with the magnetic field **H** directed at right-angles to the electric field $\vec{\mathscr{E}}$ (transverse geometry) and with both fields parallel ($\mathbf{H} \parallel \vec{\mathscr{E}}$, longitudinal geometry). In both cases the incident light reached the sample at right-angles to both electric and magnetic fields. The scattered light was suppressed by a filter which operated in the 3.5—5 μ range and was immersed directly in liquid helium. The photocurrent spectra were measured using pulsating or constant illumination. In the first case we used the method described in[5] and in the second case we employed the standard potentiometric circuit with an galvanometer and an amplifier. The measurements were carried out under constant-electric-field conditions in any one given experiment. The electric field could be varied from 0.01 to 10 V/cm. We restricted our measurements to the samples which had ohmic contacts.

The spectra recorded at 4.2° K in the absence of a magnetic field exhibited the well-known[4] oscillations of the positive photocurrent with a period of 0.026 eV, which was close to the energy of a longitudinal optical phonon ω_0. The measurements in quantizing magnetic fields showed that the sign of the transverse photocurrent changed in a certain range of frequencies of the incident radiation.

Figure 1 gives the spectral dependences of the transverse photocurrent divided by the intensity of the incident light (the light was monochromatic to with $\Delta = 1.6 \times 10^{-3}$ eV). These dependences were recorded in magnetic fields of 13, 26, 32, and 47 kOe intensity and in electric fields of 1 V/cm. The spectra in Fig. 1 were recorded for a sample with $p_0 = 2.8 \times 10^{13}$ cm^{-3} (77° K) but they were typical also of the other samples in the investigated batch.

In a magnetic field of 13 kOe (Fig. 1) the whole spectrum was positive. In stronger fields the spectrum shifted toward higher energies, the depths of the minima increased, and in a field of 18 kOe the sign of the photocurrent changed because of the excitation of photoelectrons to the two lower Landau sublevels S_0^+ and S_0^-. A further increase in the magnetic field intensity increased the depth of the minima and in a field of 32 kOe a third minimum, corresponding to the excitation of photoelectrons to the S_1^+ Landau level, crossed the zero axis and became negative.

In a field H = 47 kOe the third minimum disappeared (Fig. 1) and the depth of the remaining minima had a tendency to reach saturation with increasing magnetic field. The minimum associated with the S_1^+ level was no

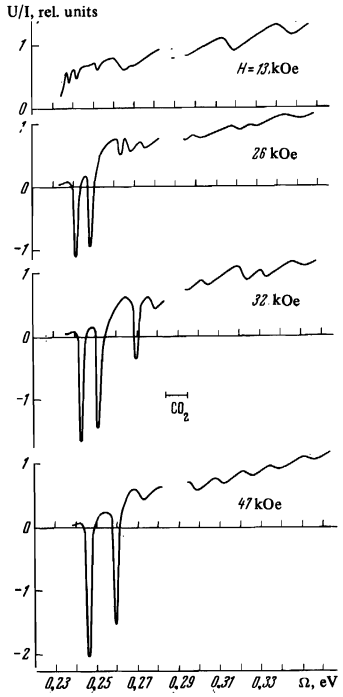

FIG. 1. Spectral dependences of the transverse photocurrent in magnetic fields H = 13, 26, 32, and 47 kOe. T = 4.2° K. The ordinate gives the photocurrent signal in relative units, calculated per unit photon flux; the abscissa is the photon energy in electron-volts.

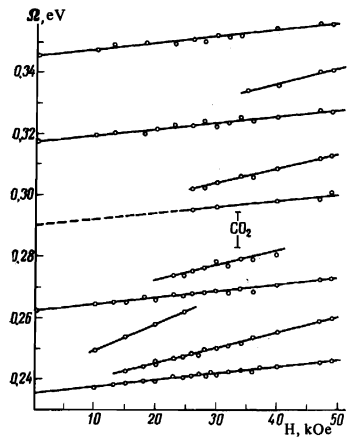

FIG. 2. Dependences of the energy positions of the minima of the transverse photocurrent on the applied magnetic field. T = 4.2° K.

longer observed because the energy of this level in fields above 32 kOe exceeded the energy of optical phonons. All the other minima were located in the positive part and they were shallower and much wider than the first two minima. These positive minima were due to the de-excitation of electrons because of the emission of optical phonons, which resulted in the transitions of these electrons to the lower Landau levels S_0^+ and S_0^- whose energies were below the optical phonon energy in all the magnetic fields used in the present study.

Figure 2 shows the dependences of the energy positions of the transverse photocurrent minima on the magnetic field. These dependences were basically similar to those reported in our earlier paper.[5] It is evident from Fig. 2 that the spectra consisted of a series of minima. The minima in each series shifted with the magnetic field in the same way as the first Landau level. The corresponding minima in different series, i.e., the

minima shifting at the same rate with the magnetic field, were separated by 0.026 eV, which was close to the energy of a longitudinal optical phonon in p-type InSb. In contrast to our earlier results,[5] some of the minima in the third and fourth series did not appear in fields below 25 kOe.

In the presence of highly monochromatic light the negative photocurrent was 10—15 times stronger than the positive current and it reached 2% of the dark current. The reversal of the sign of the transverse current due to the excitation of photoelectrons to the first Landau levels and the characteristic dependences of the positions of the photocurrent minima on the magnetic field suggested that the investigated samples of p-type InSb exhibited not only an absolute negative conductivity but also oscillations of the mobility associated with the strong nonequilibrium of the energy distribution of the photoelectrons. This was supported by measurements carried out in the longitudinal geometry and by the good agreement between the experimental and theoretical results (this point is discussed later).

In other papers[8,9] some of the present authors described spectral oscillations of the photocurrent in thick samples of n-type InSb subjected to magnetic fields. These observations were explained as being due to normal oscillations of the fundamental absorption. The samples used in these investigations were distinguished by a fairly high surface recombination velocity (in the absence of a magnetic field the spectra included a recombination peak), which led to a dependence of the photoelectron density on the absorption coefficient. This mechanism of the spectral oscillations was eliminated by carrying out a series of experiments in the longitudinal configuration (the photocurrent was measured along the magnetic field).

Since the quantization did not affect the longitudinal component of the mobility, the longitudinal photocurrent oscillations associated with the Landau levels could appear only because of oscillations of the photoelectron density. Figure 3 shows the photocurrent spectra in the transverse (curve 1) and longitudinal (curve 2) geometries for a field of 25 kOe. We can see that in the longitudinal orientation of the magnetic field there were practically no oscillations associated with the Landau levels, and, consequently, the photoelectron density did not oscillate. The absence of oscillations in the longitudinal geometry and of the recombination peak in zero magnetic

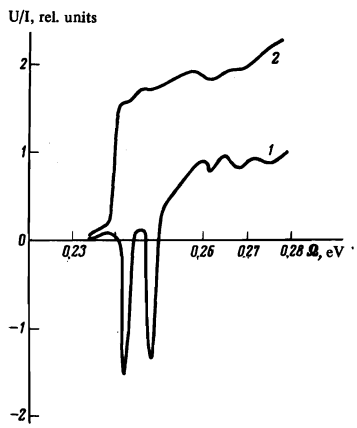

FIG. 3. Photocurrent spectra (H = 25 kOe, T = 4.2°K) in the transverse (curve 1) and longitudinal (curve 2) geometries.

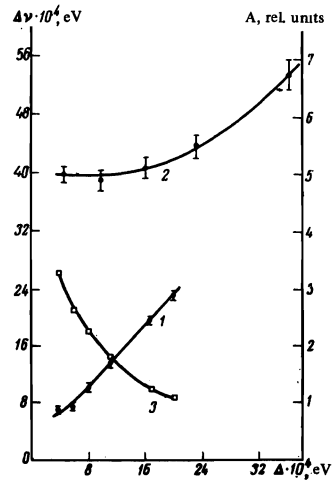

FIG. 4. Dependences of the half-width $\Delta\nu$ of the S_0^+ (curve 1) and S_{01}^+ (curve 2) minima, and of the amplitude A of the S_0^+ minimum (curve 3) on the monochromaticity Δ of the incident light. H = 39 kOe, T = 4.2°K.

field indicated that the spectral dependences of the transverse photocurrent of the investigated batch of p-type InSb samples were governed solely by the oscillations of the mobility due to the properties of the photoelectrons being far from equilibrium. We noted that the oscillations in the photocurrent spectra obtained for thick samples of n-type InSb, which were associated with the photoelectron density,[9] were independent of the orientation of the magnetic and electric fields.

The transverse photocurrent spectra (Fig. 1) depended strongly on the monochromaticity of the incident light. Figure 4 shows the dependences of the depth and spectral width (at mid-amplitude) of the first minimum S_0^+ and the width of its first "phonon replica" S_{01}^+ on the monochromaticity of the incident light, recorded in a magnetic field of 39 kOe keeping the illumination constant. Extrapolation of the linear dependence of the width S_0^+ to zero monochromaticity gave an estimate of the broadening of the lower Landau level $\Gamma \sim (1.5 \pm 1) \times 10^{-4}$ eV. The amplitude of the negative minimum S_0^+ (and also of S_0^-) in increased strongly with increasing monochromaticity (curve 3, the scale on the right). Further improvement of the monochromaticity was not possible in our experiments without reduction of the illumination intensity.

The negative photocurrent increased linearly with the intensity of the incident light (under constant monochromaticity conditions) and when the intensity exceeded 10^{15} photons · cm^{-2} · sec^{-1}, the depth of the negative minima became saturated and the photocurrent began to decrease. An increase in the electric field resulted in a linear rise of the negative photocurrent right up to $\mathscr{E} \sim 10$ V/cm. The behavior of the photocurrent oscillations in strong electric fields had several interesting properties but the discussion of these properties is outside the scope of the present paper.

3. THEORY OF PHOTOCURRENT SPECTRA IN QUANTIZING MAGNETIC FIELDS IN THE PRESENCE OF NONEQUILIBRIUM ELECTRONS

The mechanism of spectral oscillations of the photocurrent can be represented as follows (see also[5]). If the condition $\tau_R < \tau_\epsilon$ is satisfied, the photoelectrons are not in equilibrium, their mobility is negative, and it varies with the frequency of the incident light. The mobility reaches its minimum negative value when the energy of the photoelectrons approaches the Landau levels. Therefore, a minimum appears whenever the photoelectrons are excited in such a way that they acquire ener-

gies close to the lower edges of the Landau sub-bands. If the photoelectrons are created with energies in excess of the optical phonon energy, they drop down to the Landau levels "below the optical phonons" because the emission time of the optical phonons is short. This repeats the series of the minima. Such dropping of electrons to the Landau levels broadens the distribution function and reduces the depth of the phonon replicas.

We shall now give a more rigorous discussion which can be used to describe the photocurrent spectra.

A. Distribution function of photoelectrons

The state of an electron in the conduction band can be represented by the quantum numbers $\nu = (n, p_z, s)$ and p_x, where p_z and p_x are the projections of the momentum along the magnetic field and at right-angles to this field. An electron has the energy

$$\varepsilon_\nu = \varepsilon_{ns} + p_z^2/2m,$$

where $\epsilon_{ns} = \omega_c n$ for a spin oriented mainly against the field and $\epsilon_{ns} = \omega_c n + \epsilon_s$ for a spin oriented along the field; $\omega_c = eH/m$; ϵ_s is the spin splitting of the Landau levels; $n = 0, 1, \ldots$. The influence of the scattering process on the electron spectrum can be allowed for by introducing the width of the energy levels Γ_ν,[10] which is equivalent to the introduction of the relaxation time $\frac{1}{2}(\Gamma_\nu + \Gamma_{\nu'})$ for the nondiagonal components of the density matrix or for the correlation function applicable to the interaction with phonons. The energy is measured from the edge of the lowest Landau sub-band.

The transport equation for the distribution function (the diagonal component of the density matrix) is obtained by direct generalization of the corresponding equations given by Zyryanov.[11] After allowance for all possible elastic collisions, emission of optical phonons, recombination, interband transitions accompanied by the absorption of light, and weak spatial inhomogeneity (the absorption length k^{-1} is much longer than the wavelength Λ), this equation becomes

$$\sum_{\nu'} \frac{\partial^2}{\partial y^2} f_{\nu'}(y)\{D_{el}^{\nu'\nu} + D_{op}^{\nu'\nu}\} + \sum_{\nu} W_{el}^{\nu\nu'}\{f_{\nu'}(y) - f_\nu(y)\} \quad (1)$$

$$+ \sum_{\nu} \{W_{op}^{\nu'\nu} f_{\nu'}(y) - W_{op}^{\nu\nu'} f_\nu(y)\} - f_\nu(y)(\tau_R^\nu)^{-1} = -I_\nu(y),$$

where

$$D_{el}^{\nu'\nu} = \pi \sum_{\mathbf{q}} (I_{\mathbf{q}}^{\nu'\nu})^2 (C_{el}^{\mathbf{q}})^2 Y^2 \mathcal{L}(\varepsilon_\nu - \varepsilon_{\nu'}, \Gamma_\nu + \Gamma_{\nu'})$$

is the coefficient of diffusion from a state ν' to a state ν because of elastic scattering,

$$D_{op}^{\nu'\nu} = \pi \sum_{\mathbf{q}} (I_{\mathbf{q}}^{\nu'\nu})^2 (C_{op}^{\mathbf{q}})^2 Y^2 \mathcal{L}(\varepsilon_{\nu'} - \varepsilon_\nu - \omega_0, \Gamma_{\nu'} + \Gamma_\nu)$$

is the coefficient of diffusion due to emission of an optical phonon,

$$W_{el}^{\nu\nu'} = 2\pi \sum_{\mathbf{q}} (I_{\mathbf{q}}^{\nu\nu'})^2 (C_{el}^{\mathbf{q}})^2 \mathcal{L}(\varepsilon_{\nu'} - \varepsilon_\nu, \Gamma_{\nu'} + \Gamma_\nu)$$

is the probability of elastic scattering (per unit time) from the state ν to the state ν',

$$W_{op}^{\nu\nu'} = 2\pi \sum_{\mathbf{q}} (I_{\mathbf{q}}^{\nu\nu'})^2 (C_{op}^{\mathbf{q}})^2 \mathcal{L}(\varepsilon_{\nu'} - \varepsilon_\nu + \omega_0, \Gamma_{\nu'} + \Gamma_\nu)$$

is the probability of inelastic scattering by optical vibrations,

$$I_\nu(y) = \frac{8\pi^2}{\varkappa} \sum_{\mu} (\mathbf{e}\mathbf{d}_{\mu\nu})^2 \int_{-\infty}^{\infty} d\omega I(\omega) e^{-k(\omega)y} \mathcal{L}(\varepsilon_\nu - \varepsilon_\mu - \omega, \Gamma_\mu + \Gamma_\nu),$$

$y = p_x\Lambda^2$ is the coordinate of the center of the Larmor orbit in the direction of propagation of light, $Y = q_x\Lambda^2$, $C_{el}^{\mathbf{q}}$, and $C_{op}^{\mathbf{q}}$ are the Fourier components of the relevant potentials. Under steady-state illumination conditions the generation of photoelectrons can be described by

$$(I_{\mathbf{q}}^{\nu\nu'})^2 = \left| \int_{-\infty}^{\infty} dy \Phi_{\nu'}^*(y - q_x\Lambda^2) \exp(-iq_y y) \Phi_\nu(y) \right|^2 \delta(p_z, p_z' + q_z), \quad (2)$$

where μ are the states in the valence band, $I(\omega)$ is the spectral distribution of the energy flux in the incident light, κ is the refractive index of the semiconductor, $\mathbf{d}_{\mu\nu}$ is the matrix element of the dipole moment, $\mathbf{e}$ is the polarization vector of the incident light,

$$\mathcal{L}(x, \Gamma) = \frac{1}{2\pi} \frac{\Gamma}{x^2 + \Gamma^2/4}.$$

Equation (1) should be supplemented by suitable boundary conditions.

The relative order of smallness of the first term in Eq. (1) has the form

$$(k\Lambda)^2(\tau_{el}^{-1} + \tau_{op}^{-1}) / 2(\tau_R^{-1} + \tau_{op}^{-1}),$$

except in the case of scattering by impurity ions for which this ratio becomes

$$(k\Lambda)^2(\tau_{el}^{-1}E_s / \omega_c + \tau_{op}^{-1}) / 2(\tau_R^{-1} + \tau_{op}^{-1}),$$

where $E_s = k_s^2/2m$, and k_s is the reciprocal of the screening length. Usually the relaxation times are related by the following expressions:

$$\tau_{op} < \tau_{el} \ll \tau_R \text{ for } \varepsilon > \omega_0,$$
$$\tau_{op} \gg \tau_R \gg \tau_{el} \text{ for } \varepsilon < \omega_0.$$

In the first of the above two energy ranges the ratio of the diffusion term to the other terms is small $[\sim (k\Lambda)^2 \ll 1]$. In the second range the ratio depends on the lifetime and it reaches its maximum value for $p_z = 0$, i.e., under optical absorption resonance conditions (k_0 is the absorption coefficient far from resonance): $k_0^2\tau_R/2m$ (or $k_0^2\tau_R E_s/2m\omega_c$ for a long-range potential).

The relevant parameters of p-type InSb at 4.2°K are $k_0 \sim 5 \times 10^3$ cm^{-1} and $\tau_R \sim 10^{-10}$ sec. Thus, the relative magnitude of the term describing the spatial diffusion does not exceed 0.1. We can also show that even if the relative magnitude of the diffusion term is $\gtrsim 1$ it can still be ignored if the surface recombination velocity s is sufficiently low: $s\tau_R \ll L$, where L is the diffusion length. In the case of p-type InSb this condition is satisfied subject to $s \ll 10^5$ cm/sec. In all the other cases the diffusion and surface recombination lead to a dependence of the total number of photoelectrons on the frequency of the incident light and the photoelectron density oscillates with a period typical of the absorption coefficient. Such oscillations were reported earlier.[8,9] However, it is always experimentally possible to suppress the oscillations of the photoelectron density.

In view of the foregoing discussion, we shall ignore the first term in Eq. (1), i.e., we shall consider the spectral dependence of the photocurrent which is solely due to oscillations of the electron mobility. However, even in the absence of diffusion Eq. (1) remains fairly complex expression. Therefore, it is convenient to consider first a simple model which yields the positions of the resonance peaks throughout the spectral region but does not describe the general shape of the spectrum. Prior to the use of this model, we must make certain simplifying assumptions detailed below.

a) We shall assume that $\epsilon_\mu = -E_H$ and $\Gamma_\mu = \Gamma_h$, i.e., there is no dispersion in the valence band. This assumption is justified by the large difference between the masses of the heavy holes and the electrons and by the fact that the first optically active levels of the light holes in InSb practically coincide with the levels of the heavy holes,[12] so that the presence of transitions from these levels does not give rise to any additional resonance in the photocurrent spectra (Fig. 1); E_H is the forbidden band width in a magnetic field.

b) We shall assume that $\tau_{op}^{\nu} = \tau_{op}(\epsilon_\nu)$ and $\tau_R^{\nu} = \tau_R(\epsilon_\nu)$.

The assumptions a), b) allow us to seek the solution in the form $f_\nu = f(\epsilon_\nu)$. In this case the second term in Eq. (1), which describes the relaxation of the longitudinal momentum, is equal to zero.

c) We shall also ignore the broadening of the levels.

When these simplifications are introduced, Eq. (1) becomes

$$\tau_{op}^{-1}(\varepsilon_\nu + 2\omega_0) f(\varepsilon_\nu + \omega_0, y) - (\tau_{op}^{-1}(\varepsilon_\nu) + \tau_R^{-1}(\varepsilon_\nu)) f(\varepsilon_\nu, y) = -I(\varepsilon_\nu, y),$$
$$I(\varepsilon_\nu, y) = I(\varepsilon_\nu)\exp[-k(\varepsilon_\nu + E_H)y]\,8\pi^2 d^2/3\varkappa. \quad (3)$$

The solution of Eq. (3) can be found by iteration:

$$f(\varepsilon_\nu, y) = \frac{\tau_R(\varepsilon_\nu)\tau_{op}(\varepsilon_\nu)}{\tau_R(\varepsilon_\nu) + \tau_{op}(\varepsilon_\nu)} \Big\{ I(\varepsilon_\nu, y) + \sum_{k=1}^{\infty} I(\varepsilon_\nu + k\omega_0, y) \prod_{r=1}^{k} \frac{\tau_{op}^{-1}(\varepsilon_\nu + (r+1)\omega_0)}{\tau_{op}^{-1}(\varepsilon_\nu + r\omega_0) + \tau_R^{-1}(\varepsilon_\nu + r\omega_0)} \Big\}. \quad (4)$$

Each k-th term in Eq. (4) is the source of photoelectrons of energy ϵ_ν due to transitions from the state $\epsilon_\nu + k\omega_0$ accompanied by the emission of k phonons. It is interesting to consider the function

$$F(\varepsilon_\nu) = \int_0^\infty dy f(\varepsilon_\nu, y), \quad (5)$$

which governs the photocurrent reaching the whole lateral face of a thick sample. Since $\tau_{op}^{-1} \propto g(\epsilon_\nu - \omega_0)$, $k(\epsilon_\nu + E_H) \propto g(\epsilon_\nu)$, and $\tau_{op} \ll \tau_R$ for $\epsilon > \omega_0$, we finally obtain

$$F(\varepsilon_\nu) = \frac{\tau_R(\varepsilon_\nu)\tau_{op}(\varepsilon_\nu)}{\tau_R(\varepsilon_\nu) + \tau_{op}(\varepsilon_\nu)} \frac{2}{g(\varepsilon_\nu)} \sum_{k=0}^{\infty} I(\varepsilon_\nu + k\omega_0)(\varepsilon_\nu + k\omega_0 + E_H)^{-1}, \quad (6)$$

where $g(\epsilon_\nu)$ is the density of states characterized by energy ϵ_ν:[10]

$$g(\varepsilon) = \frac{eH}{4\pi^2} m^{1/2} \sum_{n,s} g(\varepsilon - \varepsilon_{ns}, \Gamma)|_{\Gamma\to 0},$$
$$g(x, \Gamma) = \left[\frac{x + (x^2 + \Gamma^2/4)^{1/2}}{x^2 + \Gamma^2/4}\right]^{1/2}. \quad (7)$$

We shall now find the required distribution without making the assumptions a)–c). We shall restrict our treatment to the case when light excites electrons to the two lower Landau sub-bands S_0^+ and S_0^-, i.e., we shall assume that the frequency of light is $\Omega < E_H + \omega_c$. In this case the distribution depends not only on the energy of photoelectrons but also on their spin, even when the operation of the source depends only on the energy. The distribution function satisfies the equation

$$[f_\downarrow(\varepsilon) - f_\uparrow(\varepsilon)](\tau_{el}^{\uparrow\downarrow}(\varepsilon))^{-1} - f_\uparrow(\varepsilon)[(\tau_R^{\uparrow}(\varepsilon))^{-1} + (\tau_{op}^{\uparrow}(\varepsilon))^{-1}] = -I_\uparrow(\varepsilon, y);$$

$$I_\uparrow(\varepsilon, y) = I_\uparrow(\varepsilon, y) + \sum_{p_z'} \{W_{op}^{\uparrow\uparrow}(\varepsilon', \varepsilon) f_\uparrow(\varepsilon') + W_{op}^{\downarrow\uparrow}(\varepsilon', \varepsilon) f_\downarrow(\varepsilon')\}, \quad (8)$$

$$(\tau_{el}^{\uparrow\downarrow}(\varepsilon))^{-1} = \sum_{p_z'} W_{el}^{\uparrow\downarrow}(\varepsilon, \varepsilon'), \quad (\tau_{op}^{\uparrow}(\varepsilon))^{-1} = \sum_{p_z'} \{W_{op}^{\uparrow\uparrow}(\varepsilon, \varepsilon') + W_{op}^{\uparrow\downarrow}(\varepsilon, \varepsilon')\}. \quad (9)$$

For the S_0^- electron we must replace $\uparrow$ with $\downarrow$. Since we are assuming that $\Gamma \ll \omega_0$, we can substitute in $\tilde{I}$ the solution of Eq. (8) with $\tilde{I} = I$ (we shall also assume that $\Omega < 2\omega_0$). The simplest distribution function is obtained in the absence of spin relaxation, i.e., if $\tau_{op}^{\uparrow\downarrow}, \tau_{el}^{\uparrow\downarrow} \gg \tau_R$ (if $I_\uparrow = I_\downarrow$, the violation of this condition does not affect the distribution for $\Omega < \omega_0$ and, consequently, it does not affect the first series of the minima in the photocurrent spectra):

$$f_\uparrow(\varepsilon) = \frac{\tau_R\tau_{op}^{\uparrow}(\varepsilon)}{\tau_R + \tau_{op}^{\uparrow}(\varepsilon)} \Big\{ I_\uparrow(\varepsilon, y) + \sum_{p_z'} W_{op}^{\uparrow\uparrow}(\varepsilon', \varepsilon) I_\uparrow(\varepsilon', y)\tau_{op}^{\uparrow}(\varepsilon') \Big\}. \quad (10)$$

It follows from Eq. (8) that

$$(\tau_{op}^{\uparrow}(\varepsilon))^{-1} = M_{op}^{\uparrow\uparrow}\, g(\varepsilon - \omega_0, \Gamma_\uparrow(\varepsilon) + \Gamma_\uparrow(\varepsilon - \omega_0)), \quad (11)$$

where M_{op} is the square of the matrix element of the transition which depends logarithmically on ϵ.

Subject to the condition a) the absorption coefficient has the following frequency dependence

$$k(\omega) = \frac{\omega e H m^{1/2}}{3\varkappa} d^2 \{g(\omega - E_H, \Gamma_\uparrow(\omega - E_H) + \Gamma_h) + g(\omega - E_H - \varepsilon_s, \Gamma_\downarrow(\omega - E_H) + \Gamma_h)\}. \quad (12)$$

Since $\Gamma \ll \omega_c$, it is sufficient to substitute $\Gamma(x)$ corresponding to $x = 0$ in all the functions $g(x, \Gamma)$.

The nature of the energy dependence $\tau_{op}(\epsilon)$ allows us to separate two groups of electrons: one with the energies $\epsilon \gtrsim \omega_0$ and the other with the energies $\epsilon < \omega_0$. The electrons in the first group are scattered mainly by the optical vibrations and their mobility is higher than that of the electrons in the second group, which are scattered elastically. However, because of the large difference between the densities, the photocurrent is not greatly affected by the first group since their contribution is small in relation to the electrons of the second group: $\sim \tau_{el}/\tau_R \ll 1$.

If Eqs. (11) and (12) are taken into account, the distribution function of the S_0^+ electrons in the second group generated by a source with a Lorentzian profile

$$I(\omega) = {}^1/_2 I_0 \Omega \mathcal{L}(\omega - \Omega, \Delta)$$

is of the form

$$F_\uparrow(\varepsilon) = I_0\tau_R \frac{4\pi^2}{eHm^{1/2}} \Big\{ \frac{\mathcal{L}(\varepsilon - \Omega + E_H, \Delta + \Gamma_0 + \Gamma_h)}{g(\varepsilon, \Gamma_0 + \Gamma_h) + g(\varepsilon - \varepsilon_s, \Gamma_0 + \Gamma_h)} + \frac{(\varepsilon + \omega_0 - \varepsilon_s)^{1/2}}{(\varepsilon + \omega_0)^{1/2} + (\varepsilon + \omega_0 - \varepsilon_s)^{1/2}} \frac{\mathcal{L}(\varepsilon - \Omega + E_H + \omega_0, \Delta + \Gamma_0 + \Gamma_h + 2\Gamma_{\omega_0})}{g(\varepsilon, \Gamma_0 + \Gamma_{\omega_0})} \Big\}, \quad (13)$$

where I_0 is the flux density of photons of frequency Ω, $\Gamma_0 = \Gamma_\uparrow(0) = \Gamma_\downarrow(\epsilon_S)$, $\Gamma_{\omega_0} = \Gamma_\uparrow(\omega_0) = \Gamma_\downarrow(\epsilon_S + \omega_0)$. For the S_0^- electrons with energies $\epsilon < \omega_0 + \epsilon_S$ the distribution differs only in the replacement, in the second term in Eq. (13), of the function $g(\epsilon, \Gamma_0 + \Gamma_{\omega_0})$ with $g(\epsilon - \epsilon_S, \Gamma_0 + \Gamma_{\omega_0})$ and by the replacement of $(\epsilon + \omega_0 - \epsilon_S)^{1/2}$ in the numerator with $(\epsilon + \omega_0)^{1/2}$. It follows from Eq. (13) that the emission of an optical phonon broadens the distribution function [second term in Eq. (13)] by an amount proportional to the reciprocal of the emission time (because $\tau_{\omega_0} \propto \tau_{op}^{-1}$).

In Eq. (1) we have ignored the energy relaxation processes (with the exception of the interaction with optical vibrations) and we have assumed that $\tau_R \ll \tau_\epsilon$. We have also ignored the possible broadening of the distribution because of plasma oscillations (these oscillations can appear at fairly high plasma frequencies

$\omega_p > \tau_R^{-1}$). To some extent all these processes can be allowed for by introducing a quantity Δ which is a measure of the broadening of the distribution provided this quantity is not related solely to the monochromaticity of the incident light.

B. Transverse conductivity of nonequilibrium photoelectrons. Form of photocurrent spectra

We shall use the well-known expression for the transverse current[13,3] and introduce the concept of the level width:

$$J_x = 2\pi e a \sum_{\nu, p_y} F_\nu \sum_{\nu', q} (I_q^{\nu\nu'})^2 X \{(C_{op}^q)^2 \mathcal{L}(\varepsilon_\nu - \varepsilon_{\nu'} - \omega_0 + e\mathcal{E}X, \Gamma_\nu + \Gamma_{\nu'}) + (C_{el}^q)^2 \mathcal{L}(\varepsilon_\nu - \varepsilon_{\nu'} + e\mathcal{E}X, \Gamma_\nu + \Gamma_{\nu'})\}, \quad (14)$$

where a is the dimension in the direction z and $X = q_y\Lambda^2$. If we assume that $e\mathcal{E}\Lambda \ll \Gamma$ (this condition is satisfied in fields $\mathcal{E} \lesssim 10$ V/cm), we can expand Eq. (14) as a series in $\mathcal{E}$ and thus find the transverse conductivity

$$\sigma_{xx} = 2e^2 a \sum_{\nu, p_y} F_\nu \frac{\partial}{\partial \varepsilon_\nu} \sum_{\nu'} (D_{op}^{\nu\nu'} + D_{el}^{\nu\nu'}). \quad (15)$$

We shall assume that the electric field is applied at right-angles to the magnetic field and to the direction of propagation of light. Moreover, we shall introduce a new quantum number p_y. The coordinate y along the direction of propagation of light will be regarded as a classical variable (the validity of this approach is discussed in detail by Zyryanov[11]).

We shall start by calculating the conductivity subject to the assumptions a)—c). For the sake of simplicity we shall assume that the source is very narrow: $\Delta \to 0$. We shall allow only for the contribution of electrons with $\epsilon < \omega_0$. Substituting F_ν from Eq. (6) into Eq. (15), we find that subsequent integration yields

$$\sigma_{xx} = -\frac{a}{2^{1/2}} I_0 \tau_R \sum_{ns, n's'} \bar{M}_{el}^{ns, n's'}(\Omega) \left(\Omega - E_H - \omega_0 \left[\frac{\Omega - E_H}{\omega_0}\right] - \varepsilon_{n's'}\right)^{-3/2}$$

$$\varepsilon_{ns}, \varepsilon_{n's'} \leqslant \Omega - E_H - \left[\frac{\Omega - E_H}{\omega_0}\right]\omega_0. \quad (16)$$

Here,

$$\bar{M}_{el}^{ns, n's'}(\Omega) = \frac{m^{1/2}}{4H^2} \sum_q (C_{el}^q)^2 (I_q^{ns, n's'}(\Omega))^2 q_y^2$$

is a smooth positive function of Ω (in the case of a short-range potential this function is independent of Ω) and [x] is the integral part of a number.

In this way we obtain a negative oscillating conductivity with minima at source frequencies satisfying the condition

$$\Omega = E_H + \varepsilon_{n's'} + \omega_0 \left[\frac{\Omega - E_H}{\omega_0}\right], \quad \varepsilon_{n's'} < \omega_0. \quad (17)$$

In the earlier paper[5] we pointed out that this condition yields series of minima separated from one another by the optical phonon frequency.

We should note that, in contrast to the logarithmic divergence in the equilibrium case,[10] the transverse conductivity of strongly nonequilibrium electrons diverges as $\delta^{-3/2}$, where δ is a combination of the level width and the spectral width. The nature of the photocurrent spectrum is very sensitive to the broadening of the Landau levels. This is confirmed by the numerical calculations given below.

The following expression can be obtained for the total photoconductivity which includes the contribution of photoholes and of the Hall emf:

$$\Delta\sigma = (\sigma_{xx} + \sigma_{xx}^h)\left[1 - \left(\frac{\sigma_{xy}}{\sigma_{xx}}\right)_d^2\right] + 2(\sigma_{xy} + \sigma_{xy}^h)\left(\frac{\sigma_{xy}}{\sigma_{xx}}\right)_d, \quad (18)$$

where the index h refers to the photoholes and the subscript d refers to holes in darkness. Substituting in Eq. (18) the expression for σ_{xx} with F_ν from Eq. (13), we obtain

$$\Delta\sigma = C_1 \int_0^\infty \frac{dx}{x^{1/2}} \left\{\frac{\partial}{\partial x} g(x, 2\Gamma_0) \left[\frac{\mathcal{L}(x - \Omega + E_H, \Gamma_0 + \Gamma_h + \Delta)}{g(x, \Gamma_0 + \Gamma_h) + g(x - \varepsilon_s, \Gamma_0 + \Gamma_h)} + \frac{\mathcal{L}(x - \Omega + E_H + \varepsilon_s, \Gamma_0 + \Gamma_h + \Delta)}{g(x, \Gamma_0 + \Gamma_h) + g(x + \varepsilon_s, \Gamma_0 + \Gamma_h)} + \frac{(x - \varepsilon_s + \omega_0)^{1/2}}{(x + \omega_0)^{1/2} + (x - \varepsilon_s + \omega_0)^{1/2}} \times \frac{\mathcal{L}(x - \Omega + E_H + \omega_0, \Delta + \Gamma_0 + \Gamma_h + 2\Gamma_{\omega_0})}{g(x, \Gamma_0 + \Gamma_{\omega_0})} + \frac{(x + \omega_0 + \varepsilon_s)^{1/2}}{(x + \omega_0)^{1/2} + (x + \varepsilon_s + \omega_0)^{1/2}} \frac{\mathcal{L}(x - \Omega + E_H + \omega_0 + \varepsilon_s, \Delta + \Gamma_0 + \Gamma_h + 2\Gamma_{\omega_0})}{g(x, \Gamma_0 + \Gamma_{\omega_0})}\right] + C \frac{\mathcal{L}(x - \Omega + E_H, \Gamma_0 + \Gamma_h + \Delta)}{g(x, \Gamma_0 + \Gamma_h)}\right\}; \quad (19)$$

$$\Delta n = \Delta p \frac{\tau_R}{\tau_h} = I_0 \tau_R \sqrt{2} \int_0^\infty \frac{dx}{x^{1/2}} \frac{\mathcal{L}(x - \Omega + E_H, \Gamma_0 + \Gamma_h + \Delta)}{g(x, \Gamma_0 + \Gamma_h)}$$

$$C_1 = a\sqrt{2} I_0 \tau_R \bar{M}_{el}^{00} \left[1 - \left(\frac{\sigma_{xy}}{\sigma_{xx}}\right)_d^2\right],$$

$$C = e(\bar{M}_{el}^{00})^{-1} \left\{\frac{\mu_{xx}^h \tau_h}{\tau_R} + \frac{2(\mu_{xy} + \mu_{xy}^h \tau_h/\tau_R)}{1 - (\sigma_{xy}/\sigma_{xx})_d^2} \left(\frac{\sigma_{xy}}{\sigma_{xx}}\right)_d\right\}.$$

Here, Δn is the photocarrier density in a sample with one illuminated surface; it is assumed that the photohole conductivity is governed by the expression $\mu\Delta p$, where μ is the frequency-independent mobility. This assumption is justified by the fact that photoholes are created with very low energies and, moreover, they become rapidly thermalized so that their lifetime τ_h is quite long.

The computer-calculated spectrum (19) was compared with the experimental curve obtained in H = 39 kOe. Different values of the parameters in Eq. (19) were used in this comparison. Figure 5 shows the theoretical and experimental dependences of the photocurrent on the frequency of the incident light for $\Delta = 1.8 \times 10^{-3}$ eV, $E_H = 0.244$ eV, $\omega_0 = 0.022$ eV, $\epsilon_s = 0.0092$ eV, $\Gamma_h = 0.4 \times 10^{11}$ sec^{-1}, $\Gamma_0 = 1.85 \times 10^{11}$ sec^{-1}, $\Gamma_{\omega_0} = 4.6 \times 10^{12}$ sec^{-1}, $C = 0.13$ meV$^{-3/2}$. Numerical calculations showed that the form of the spectrum (particularly the depth, posi-

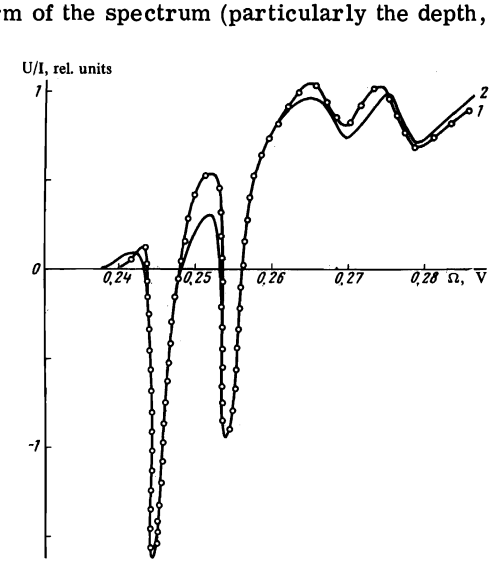

FIG. 5. Spectra of the transverse current in a magnetic field H = 39 kOe, T = 4.2°K: 1—theoretical spectrum; 2—experimental results.

tions, and widths of the minima) is a function of each of these parameters. A change in E_H shifts the whole spectral curve along the energy axis. The spin splitting and the optical phonon frequency determine the separation between the minima. The spectrum becomes more symmetric, i.e., the positive peaks in front of the first and second minima increases, when the parameter Γ_h is increased. The value Γ_0 determines primarily the depth of the first and second minima. The width and the depth of "phonon replicas" depend on Γ_{ω_0}. If C is increased, the whole spectrum shifts to the positive region, in ag agreement with the experimental results reported in[5].

4. CONCLUSIONS

The reported experimental results on the negative photocurrent and the good agreement between the experimental and theoretical spectra demonstrate that the absolute negative conductivity of the photoelectrons in p-type InSb appears at helium temperatures if quantizing magnetic fields are applied. As pointed out in the Introduction, an absolute negative conductivity can alter the sign of the total current. The appearance of this conductivity is hindered by the positive background due to the equilibrium majority carriers (the dark current). In our experiments the negative signal reached 2% of the dark conductivity. The absolute negative conductivity may become total in samples with a high degree of compensation illuminated with strong light, such as a laser beam with good monochromaticity. It would also be interesting to determine the spectral distribution of the photocurrent in other semiconductors subjected to quantizing magnetic fields. These investigations should be carried out on materials for which a nonequilibrium energy distribution of electrons can be expected. It is worth noting the experimental observations of the oscillations of the photocurrent in GaSb[6] which may be attributed to an absolute negative conductivity in quantizing magnetic fields.

Investigations of the spectral oscillations of the photocurrent due to strongly nonequilibrium photoelectrons are also of intrinsic interest because the sensitivity of such oscillations to the broadening of the Landau levels, the emission time and the frequency of the optical phonons, and the energy band structure in quantizing magnetic fields should make it possible to determine these characteristics from the experimental curves.

Apart from oscillations of the photocurrent, the absolute negative photoconductivity can lead to several other effects which may be associated with the instability of the spatial and energy distributions of nonequilibrium electrons present in sufficiently high densities.[14]

The authors are grateful to P. L. Nevskiĭ for his great help in the computer calculations.

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Translated by A. Tybulewicz
25