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Joint scientific session of the Physical Sciences Division of the Russian Academy of Sciences and the Joint Physical Society of the Russian Federation "Absolute Negative Conductivity" (27 October 2004)

A joint scientific session of the Physical Sciences Division of the Russian Academy of Sciences (RAS) and the Joint Physical Society of the Russian Federation dedicated to "Absolute Negative Conductivity" was held in the Conference Hall of the P N Lebedev Physics Institute, RAS, on 27 October 2004. The following reports were presented at the session:

(1) **Elesin V F** (Moscow Engineering Physics Institute (State University), Moscow) "Absolute negative conductivity phenomena in non-equilibrium three-dimensional semiconductors";

(2) **Gantmakher V F, Zverev V N** (Institute of Solid State Physics, RAS, Chernogolovka, Moscow Region) "Magneto-impurity resonances as indicators of an inverse photoelectron distribution function in semiconductors";

(3) **Ryzhiĭ V I** (University of Aizu, Aizu-Wakamatsu, Japan) "Microwave-induced negative conductivity and zero-resistance states in two-dimensional electronic systems: history and current status";

(4) **Dorozhkin S I** (Institute of Solid State Physics, RAS, Chernogolovka, Moscow Region) "Millimeter-wave response in the magnetoconductivity of highly perfect two-dimensional electron systems".

An abridge version of the papers is given below.

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Absolute negative conductivity phenomena in non-equilibrium three-dimensional semiconductors

V F Elesin

1. Introduction

This talk presents an analytical review of research on absolute negative conductivity (ANC) in non-equilibrium 3D semiconductors and covers ANC mechanisms operating in semiconductors with strong electron – optical-phonon coupling and absolute negative transverse conductivity in a quantizing magnetic field.

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At the end of 2002, teams of US, German, and Israeli scientists discovered a zero-resistance state in a 2D electron gas in a GaAs/AlGaAs semiconductor heterostructure simultaneously exposed to a constant magnetic field and electromagnetic radiation [1]. First interpreted by the authors in terms of superconductivity, the discovery was afterwards found to be closely related to absolute negative conductivity, a phenomenon in which a direct current flows opposite to the externally applied constant electric field. Despite a large body of publications on both theoretical and experimental aspects of the subject (see Refs [2, 3] for a review), the nature of the phenomenon is not yet completely understood.

ANC effects, interestingly, were predicted more than 30 years ago for 3D [4, 5] and 2D [6] systems and subsequently given a detailed theoretical and experimental study [7 – 11] for instability, domain formation, etc. It is therefore of interest here to briefly look at ANC mechanisms in non-equilibrium semiconductors. We will restrict our discussion to three-dimensional systems.

A realistic ANC mechanism seems to have been first suggested in Ref. [4] in 1966; in the same year, a detailed theoretical analysis of this mechanism was performed [7]. According to Ref. [4], the ANC mechanism involves a non-uniform electron energy distribution and employs the threshold nature of electron interaction with (dispersionless) longitudinal optical phonons.

Early in 1968 it was predicted [5, 8] that for a semiconductor placed in a quantizing magnetic field, the transverse current is negative if the electron energy distribution is non-equilibrium (ANC in a magnetic field). This effect was later discovered experimentally and studied in detail in p-InSb [9 – 11] and p-Ge [12].

It should be noted that still earlier Krömer [13, 14] suggested that a negative electron mass might lead to negative conductivity in semiconductors — without, however, giving any formulas or calculations to support the idea. In 1960, conditions for a negative mass cyclotron resonance to be able to emit radiation were formulated by Kagan [15]

2. ANC effect for non-equilibrium electrons coupled to optical phonons

The study of strongly non-equilibrium phenomena in semiconductors was triggered by experimental work [16, 17] on the observation of oscillations in photoconductivity spectra. On the one hand, the oscillation period was found to be equal to the frequency of an optical phonon, but the most surprising finding was that photonconductivity was dependent at all on the frequency of the exciting light. At the time it was the accepted view that the energy distribution of non-equilibrium photoelectrons differs very little from the quasi-equilibrium distribution [18]. What the frequency dependence suggested was that the conduction band electron lifetime τ_e is comparable to the energy relaxation times due to acoustic phonon scattering (τ_{ac}) and to electron – electron collisions (τ_{ee}). In the limit $\tau_e \ll \tau_{ac}, \tau_{ee}$ the electron distribution function would be δ-like, corresponding to a strongly non-equilibrium state, and it turned out that such a situation occurs, for example, in p-type indium antimonide [17], in which $\tau_e = 10^{-10}$ s, $\tau_{ac} = 10^{-7}$ s.

In Ref. [4] it was shown that a strongly non-equilibrium state may exhibit absolute negative conductivity. Following Refs [4, 7], let us consider a semiconductor at low temperatures $kT \ll \hbar\omega_0$, where ω_0 is the optical phonon frequency (from now on we take $\hbar = 1$). We assume that equilibrium electrons have a lower concentration compared to the conduction band electrons of energy ε_0 created by an external monochromatic source of intensity I with distribution $g(\Omega)$ (Fig. 1).

The behavior of electrons in the presence of an external electric field $\mathbf{E}$ is described by the usual kinetic equation for the electron distribution function $f(\mathbf{p})$,

$$-e\mathbf{E}\,\frac{\partial f(\mathbf{p})}{\partial \mathbf{p}} + S_{\mathbf{p}}^{-} f(\mathbf{p}) = S_{\mathbf{p}}^{+}\{f\} \tag{1}$$

in combination with the current expression

$$\mathbf{J} = -e\int \frac{2\,\mathrm{d}^3 p}{(2\pi)^3}\,\mathbf{v} f(\mathbf{p})\,. \tag{2}$$

The electron drift operator $S_{\mathbf{p}}^{-}$ and integral source term $S_{\mathbf{p}}^{+}$ include interaction with optical phonons and impurities (with a characteristic time τ_{imp}), as well as electron recombination and creation processes. In particular, the 'drift probability' due to the emission of optical phonons is expressed as

$$\frac{1}{\tau_{op}(\varepsilon_p)} = A\sqrt{\varepsilon_p - \omega_0}\,\Theta(\varepsilon_p - \omega_0)\,,$$

$$\Theta(x) = \begin{cases} 1, & x > 0 \\ 0, & x < 0\,. \end{cases} \tag{3}$$

Let us assume that $\varepsilon_0 < \omega_0$. The electrons are then distributed in phase space in a thin spherical layer around the sphere $p^2/2m = \varepsilon_0 < \omega_0$ (Fig. 2a). Because $\varepsilon_0 < \omega_0$, it follows that, according to the conservation of energy, the electrons cannot emit phonons [see Eqn (3)]. Turning on an electric field shifts the sphere in momentum by $e\mathbf{E}\tau$ as shown in Fig. 2b. The electrons that increase their energy to ω_0 (on the right semi-sphere) emit optical phonons, losing almost all

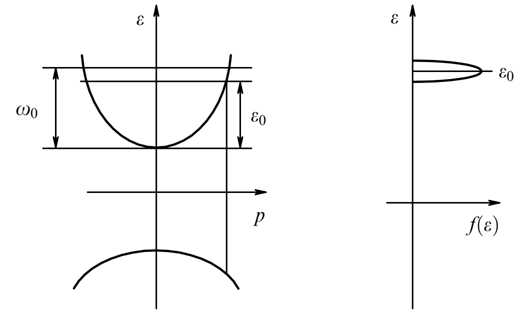

Figure 1.

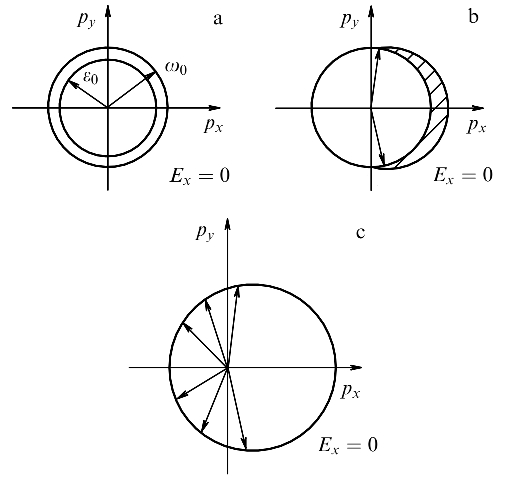

Figure 2.

of their momentum and energy in doing so (Fig. 2c). As a result, a directional current of high-momentum ($p \sim \sqrt{m\omega_0}$) electrons not coupled to phonons arises in the direction opposite to the field — contributing a negative amount $\mathbf{J}_-$ to the total current. Clearly, there is also a positive contribution $\mathbf{J}_+$, due to the momentum imparted by the field. It was shown in Refs [4, 7] that under certain conditions, the negative contribution may predominate, thus leading to ANC.

So the ANC effect is due to the fact that the field operates in a twofold manner, either directly imparting momentum to the electrons or controlling dissipative interaction processes (i.e., involving optical phonons). This duality is clearly seen if Eqn (2) is rewritten, following Ref. [7], as

$$\mathbf{J} = \mathbf{J}_1 + \mathbf{J}_2\,,$$

$$\mathbf{J}_1 = \mathbf{E}\,\frac{e^2}{m}\int \frac{2\,\mathrm{d}^3 p}{(2\pi)^3}\,S_p^{+}\{f_0\}\int_0^{\infty} \mathrm{d}t\,t\exp\left[-\int_0^t \frac{\mathrm{d}t'}{\tau(\varepsilon_p - eEt')}\right], \tag{4}$$

$$\mathbf{J}_2 = -\frac{e}{m}\int \frac{2\,\mathrm{d}^3 p}{(2\pi)^3}\,\mathbf{p} S_p^{+}\{f_0\}\int_0^{\infty} \mathrm{d}t\exp\left[-\int_0^t \frac{\mathrm{d}t'}{\tau(\varepsilon_p - eEt')}\right], \tag{5}$$

where

$$f_0(\varepsilon_p) = \frac{1}{\sqrt{2m\varepsilon_p}}\int \frac{\mathrm{d}^3 p'}{4\pi}\,f(p')\,\delta(\varepsilon_p - \varepsilon_{p'}) \tag{6}$$

is the symmetric part of the distribution function [19] and τ is the total relaxation time.

If the time τ is assumed to be independent of the energy, the current $J_2 = 0$ and $J_1 = e^2 E n\tau/m$ (n is the electron concentration). So the current J_1 is due to the momentum imparted by the field, and its magnitude is equal to that of J_+. The way the current J_+ increases with the field is shown in Fig. 3.

The current J_2 is of a totally different nature. Here the electric field acts as a factor controlling the way in which the energy from an external source is distributed between the electrons and phonons. For our case, $J_2 \simeq J_-$.

So if $\tau = $ const (in which case the field has no effect on relaxation), the current is always positive and equal to $J = e^2 E\tau n/m$ independent of the shape of the distribution

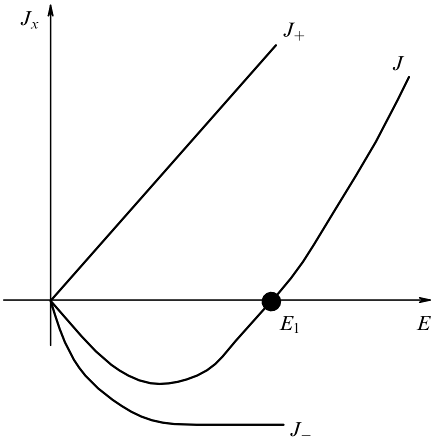

Figure 3.

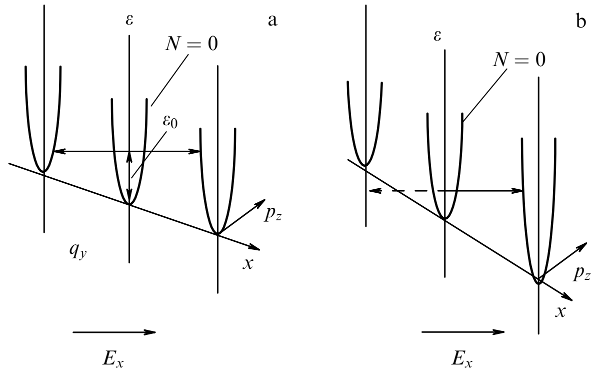

Figure 4.

function $f_0(\varepsilon)$ (in particular, also for $\partial f_0/\partial\varepsilon > 0$ in a certain energy range). Specifically, the current J_- is calculated to be

$$J_- \simeq -\frac{e^2 E\tau}{m}\, n_1 \left(\frac{\varepsilon_0}{\varepsilon_0 - \omega_0}\right)^{1/2} \xi_1\,, \qquad \xi_1 \simeq 1 \tag{7}$$

for a weak field $eE\tau\sqrt{\varepsilon_0/m} < \varepsilon_0 - \omega_0$, and

$$J_- \simeq -\frac{e^2 E\tau}{m}\, n_1 \left[\frac{\sqrt{m\varepsilon_0}}{eE\tau}\right]^{1/2} \xi_2\,, \qquad \xi_2 \sim 1 \tag{8}$$

for a strong field $eE\tau\sqrt{\varepsilon_0/m} > \varepsilon_0 - \omega_0$. Here n_1 is the concentration of electrons created by light. Upon further increase in the field the current J_- saturates because those electrons that obtained the energy from the field lose the energy. The resulting current behaves as shown in Fig. 3. At $E = E_1$ it vanishes and then becomes positive.

It is natural to expect that an ANC state can be unstable [20]: at $\sigma < 0$ the Maxwell time $\tau_M = 4\pi\sigma/\varepsilon$ becomes negative, with the result that charge fluctuations will no longer be suppressed and the charge density will increase instead.

The ANC stability study of Ref. [7] has identified the formation and stability conditions of the non-uniform state. In particular, it was shown that at, for example, $E = E_1$, $J(E_1) = 0$, a state (domain) with a spontaneous electric field may form.

3. ANC effect in a transverse quantizing magnetic field

In early 1968 it was shown [5, 8] that a monoenergetic electron distribution may lead to ANC in a transverse magnetic field H — an effect in which a current perpendicular to the magnetic field flows in the direction opposite to the applied electric field. This effect is due to the electric field affecting the impurity scattering probability, along with the peculiarities of the electron density of states in magnetic field.

Let us assume the same conditions as in Section 2 and consider a semiconductor placed in a magnetic field $\mathbf{H} \parallel \mathbf{z}$. We will assume that electrons are produced with energy ε_0 less than ω_0, that the lifetime τ_e is much shorter than τ_{ac}, τ_{ee}, and that momentum scattering is due to impurities. Note that due to one-dimensionality, the electron – electron relaxation time in the lowest Landau band is infinite [21].

It is known that in a strong magnetic field electrons are localized in a plane perpendicular to the field, close to the centers of the Larmor orbits. The electric field, while not delocalizing the electrons, changes their energy along the x-axis (Fig. 4). Scattering by impurities causes the electrons to hop a distance of about $L = (\hbar c/eH)^{1/2}$ between the orbit centers. Below, it will be rigorously proved that the hopping probability is proportional to the densities of initial and final states.

The question now arises as to which way an electron is more likely to hop, along the field (in which case the electron takes energy away from the field) or against the field. From Fig. 4, it is clear that for (elastic) impurity scattering, an electron that hops along the field comes to a state with a lower density of states (compared to a hop against the field) if the energy $\varepsilon_0 > (\Omega/2) + LE$. That is, non-equilibrium electrons with an energy $\varepsilon > \varepsilon_0$ exhibit the ANC effect.

We proceed now to a rigorous analysis along the lines of Ref. [8]. The transverse current density J_x due to the electric field $E_x \equiv E$ can be expressed in the form

$$J_x = \frac{2}{(2\pi)^2} \sum_{N, M, q_x, q_y} \int \mathrm{d}\varepsilon \left[f(\varepsilon) \times \frac{q_y |V_{\mathbf{q}, N, M}|^2}{(\varepsilon - N - 1/2)^{1/2} (\varepsilon - M - 1/2 + q_y F)^{1/2}} \right]. \tag{9}$$

Here, the quantum numbers N, M (0, 1, 2, ...) label the Landau bands, q_y is the separation between the Larmor orbit centers, $f(\varepsilon)$ is the electron distribution function, and $V_{\mathbf{q}, N, M}$ are the impurity scattering matrix elements; the units for length, momentum, energy, and force F are taken to be $L = (\hbar c/eH)^{1/2}$, $\hbar/L$, $\hbar\Omega$, and $\hbar\Omega/L$, respectively.

Expression (9), obtained by extending the method of Adams and Holstein [22], holds for strong magnetic fields and any electric fields, its easy-to-grasp form clearly demonstrating the mechanism of transverse conduction in a quantizing magnetic field. A point to note from Eqn (9) is that the current is proportional to the scattering probability, the initial and final densities of states, and the distribution function. The electric field enters the density of the final states, producing a probability difference between the energy-gain ($q_y > 0$) and energy-loss ($q_y < 0$) transitions — and hence creating a transverse electric current.

If it is assumed that $f(\varepsilon) = \delta(\varepsilon - \varepsilon_0)\tau_e$, $N = M = 0$ (which leaves us with only the lowest Landau band), then the current, Eqn (9), takes the form

$$J_x = \frac{2\tau_e}{(2\pi)^2} \sum_{q_y} \frac{q_y |V_{q,0,0}|^2}{(\varepsilon_0 - 1/2)^{1/2} (\varepsilon_0 - 1/2 + q_y F)^{1/2}}\,. \tag{10}$$

From (10) and Fig. 4a it follows that for energies $\varepsilon_0 > 1/2 + |q_y|F$ a transition with a loss of energy in the electric field to a state with a higher density of states is more likely to occur. Thus, in a wide energy range $\varepsilon_0 > 1/2 + LF$ an electron makes a negative contribution to the current, and only in a narrow range $1/2 < \varepsilon_0 < 1/2 + LF$ (Fig. 4b) does it contribute positively because energy-loss transitions are impossible. Hence, if the energy of non-equilibrium electrons exceeds $1/2 + LF$, the ANC effect takes place.

We next present the results of the calculation of the current for two extreme cases. If the electric field is small, $F \ll \varepsilon_0 - 1/2$, Eqn (10) yields

$$J_x = -\frac{\beta F \tau_e}{(\varepsilon_0 - 1/2)^2}\,, \tag{11}$$

or for the transverse conductivity in dimensional form,

$$\sigma_{xx} = -\sigma_0 \left(\frac{3\Omega}{32\varepsilon_0}\right) \frac{\Omega^2}{(\varepsilon_0 - \Omega/2)^2} < 0\,, \tag{12}$$

where σ_0 is the 'classical photoconductivity' in a magnetic field and β is a numerical factor. Note that as $\varepsilon_0 \to \Omega/2$, the ANC greatly exceeds σ_0. In the opposite limit, $F \gg \varepsilon_0 - 1/2$, the current

$$J_x = \frac{\gamma \tau_e}{\sqrt{F}\sqrt{\varepsilon_0 - 1/2}} > 0\,, \quad \gamma \sim 1\,, \tag{13}$$

becomes positive.

Figure 3 shows the qualitative behavior of $J_x(E)$. Note that the field dependence of J_x is similar to that for the current of electrons interacting with optical phonons (see Section 2), so that the instability analysis of Ref. [7] is valid here as well. In particular, for $E_x = E_1$, $J_x(E_1) = 0$, and we have a state with $\sigma_{xx} = 0$. As is known, in a strong magnetic field σ_{xx} is proportional to the resistance R_{xx}, which also turns out to be zero, $R_{xx}(E_1) = 0$ — which is exactly what confused the authors of Ref. [1].

It should be noted that the above result does not depend on specifically how the non-equilibrium state was created — except for the point that the non-equilibrium electrons should be distributed in the energy range $\varepsilon > 1/2 + LF$.

4. Experiments and conclusions

The ANC effect, as predicted in Ref. [5], was in due course discovered in p-InSb in 1970 [9] and then investigated in detail in Refs [10, 11]. A typical variation of σ_{xx} with monochromator frequency (i.e., energy ε_0) is shown in Fig. 5 for $T = 4.2$ K, $H = 39$ kOe [10]. The two sharp minima on the graph correspond to the resonances at the first ($\varepsilon_0 = \Omega/2$) and second ($\varepsilon_0 = 3\Omega/2$) Landau levels. After that the electron energy becomes equal to the energy of an optical phonon, whose emission produces a broadening of the resonance and leads to a positive conductivity. A study of ANC in Ref. [11], over a wide range of magnetic fields (up to 150 kOe), provided, in particular, a unique insight into the parameters of the semiconductor.

The experiments of Refs [9 – 12] proved the existence of the ANC effect in photoelectrons and helped to identify the mechanism of transverse conduction, but the low intensity of the monochromator and the relatively high contribution of equilibrium holes prevented the observation of ANC in full. In the past thirty years, remarkable progress has been made in technology and experimentation, as witnessed by the discovery of the ANC phenomenon in two-dimensional systems. It seems reasonable to concentrate on three-dimensional systems in further work as they exhibit a much stronger effect than two-dimensional ones.

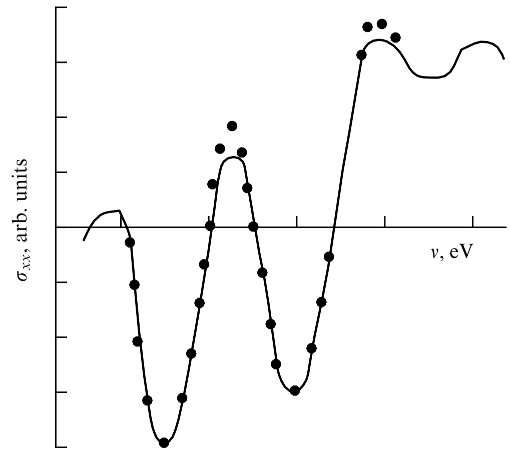

Figure 5.

The reason is that in two-dimensional systems ANC is due to photon absorption occurring simultaneously with scattering by an impurity. This is a weak process similar to indirect transitions in semiconductors, which must involve phonons for a photon to be absorbed. It should be noted that, given that a complete theory for 2D heterostructures is still lacking, the experimentally supported earlier theory can be useful in providing a thorough interpretation for the phenomena discovered in Ref. [1]

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Magnetoimpurity resonances as indicators of an inverse photoelectron distribution function in semiconductors

V F Gantmakher, V N Zverev

The mechanism of absolute negative conductivity, due to the specific way in which electrons behave in crossed electric and quantizing magnetic fields, was predicted by Elesin in 1968 [1] and observed experimentally by Aleksandrov et al. [2]. According to Ref. [1], because the magnetic subband density of states $g(\varepsilon) \sim \varepsilon^{-1/2}$ is a decreasing function of energy, an electron which is far enough above the bottom of a subband tends to gain potential energy at the expense of its kinetic energy when elastically (or quasi-elastically) scattered in collisions. The energy conservation law for such collisions is

$$eE\Delta X + \Delta\varepsilon \pm u = 0\,, \tag{1}$$

where the first two terms are the respective changes in the electron potential and kinetic energies, ΔX is the shift in the position of the Larmor orbit center, and u is the energy of the absorbed (emitted) acoustic phonon. (The term u in Eqn (1) is added for phonon scattering, for which the quasi-elasticity condition is expressed by the inequality $\varepsilon \gg u$.)

For electrons with energy $\varepsilon > eE\Lambda$ the average value of $eE\Delta X$ is positive. For an equilibrium distribution function, this negative contribution to the conduction is compensated by the positive contribution from the electrons in the energy range $\varepsilon \leqslant eE\Lambda$ close to the subband bottom. But if such near-bottom electrons are too few for some reason, then in the presence of a strong magnetic field and a perpendicular electric field the total conductivity of this group of carriers may prove negative.

In the experiment in Ref. [2], non-equilibrium electrons were produced by interband monochromatic photoexcitation in the direct-band-gap semiconductor p-InSb at helium temperatures. Because the lifetime of the photoelectrons was much shorter than the characteristic energy relaxation time, the electron distribution function turned out to be quite far from equilibrium and was determined by the line shape of the light source used. Therefore, any time electrons were created close to the edges of the magnetic subbands, the photocurrent spectrum exhibited a negative minimum, due to the absolute negative conductivity effect.

A necessary but clearly not sufficient condition for absolute negative conductivity to occur is that there be a region in which the distribution function has a positive derivative, $\partial f/\partial\varepsilon > 0$. This point seems to have been first made in the theoretical work of Ref. [3], in which the cyclotron resonance response in the conductivity of semiconductors with negative effective masses was investigated. For low-frequency transport, in crossed electric and magnetic fields the easiest way to show this is by means of a model which describes exchange processes in an electron system in terms of electron diffusion in energy space [4, 5]. In this model, the energy obtained from the electric field, $(\partial U/\partial t)_E$, is given by

$$\left(\frac{\partial U}{\partial t}\right)_E = -\int_0^\infty D_\varepsilon \frac{\partial f}{\partial \varepsilon}\, g\, \mathrm{d}\varepsilon\,, \tag{2}$$

where D_ε, the diffusion coefficient along the ε-axis, is determined by how far the center of the Larmor orbit shifts in collisions [5]. From Eqn (2) it is seen that the inequality $(\partial U/\partial t)_E < 0$ can only be satisfied if there is a region where $\partial f/\partial\varepsilon > 0$. In addition, it can be seen that for a given function $f(\varepsilon)$ integral (2) can be positive for one function D_ε (i.e., for one type of scatterers) and negative for another.

Absolute negative conductivity mechanisms for inelastic processes in two-dimensional systems were analyzed theoretically by Ryzhiĭ [6], whose predictions were to wait 30 years before being experimentally realized. Relatively recently, zero-resistance states were discovered in samples of 2D high electron mobility GaAs/AlGaAs heterostructures subjected to strong microwave radiation [7]. In describing this phenomenon, current models [8] are all to a greater or lesser extent based on the mechanism of absolute negative conductivity — an effect due to a strong non-equilibrium which microwave radiation produces in the electron distribution function [1, 6].

Experiments on heterostructures renewed interest in the subject, and it seems therefore appropriate here to revisit our old experiments on low-temperature magnetotransport in photoexcited germanium doped with shallow acceptor impurities [9]. As part of that work, a study of magnetoimpurity oscillations was carried out, which showed that conductivity in crossed electric and quantizing magnetic fields has a negative part, due to the distortion of the photoelectron distribution function. The magnetoimpurity oscillations look like a sequence of photocurrent peaks periodic in an inverse magnetic field (Fig. 1), which are due to inelastic

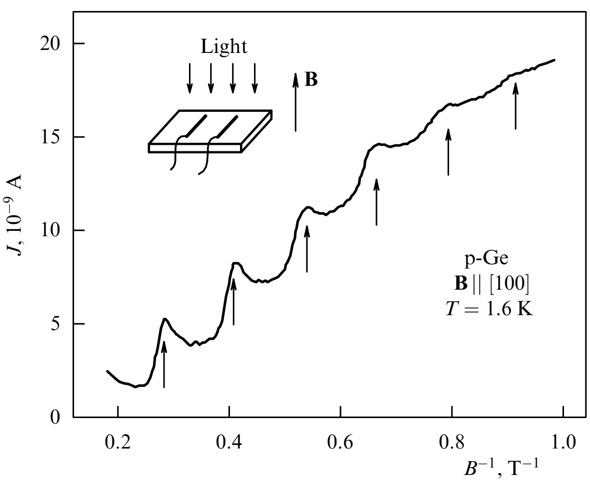

Figure 1. Magnetoimpurity oscillations of photoconductivity in a sample of gallium-doped p-Ge. $N_{\mathrm{Ga}} = 2\times10^{14}$ cm^{-3}, the pulling field $E = 3.5$ V cm^{-1}, interband generation rate $G = 10^{17}$ cm^{-3} s^{-1}, light wavelength $\lambda = 0.63$ μm.