ANISOTROPIC PHOTOELECTRIC EFFECTS

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Photoelectric effects, caused by polarized radiation in media whose optical properties are isotropic, are investigated. Under conditions for the observation of photoconductivity, side by side with it a photo-emf arises in the direction perpendicular to the constant field. Mechanisms for the origin of a transverse photo-emf and the dependence of the photoconductivity on polarization of the radiation in many-valley semiconductors or in an isotropic plasma are considered.

1. PHOTOCONDUCTIVITY TENSOR AND ANISOTROPIC PHOTOELECTRIC EFFECTS

One can regard photoconductivity as a nonlinear electrodynamical effect of the third order.[1] Actually the photocurrent is proportional to the first power of the constant (drawing) field and to the second power of the amplitude of the variable (microwave or optical) field. Therefore, in a homogeneous medium the time-independent photocurrent can be represented in the form

$$j_i^{\mathrm{ph}} = \overset{\mathrm{ph}}{\sigma}_{ijkl} E_j F_k F_l. \qquad (1)$$

Here E_i and F_i are, respectively, the components of the constant field in the sample and the amplitudes of the variable field of frequency ω. One can call the tensor $\sigma^{\mathrm{ph}}_{ijkl}$ the photoconductivity tensor. From the definition (1) it follows that one can regard it as symmetric with respect to permutation of the last two indices.

We shall investigate the photoelectric effects which follow from expression (1) for photocurrents in those media whose optical properties and static electrical conductivity are isotropic, that is, in isotropic media and cubic crystals. In these media the following components of the photoconductivity tensor do not vanish: $\sigma^{\mathrm{ph}}_{iiii} = \gamma_{11}$, $\sigma^{\mathrm{ph}}_{iijj} = \gamma_{12}$, and $\sigma^{\mathrm{ph}}_{ijij} = \gamma_{44}$ ($i \neq j$). The other components are equal to zero. In addition, in an isotropic medium one has

$$\gamma_{44} = (\gamma_{11} - \gamma_{12})/2. \qquad (2)$$

We shall show that under the conditions for observation of photoconductivity, a transverse field (perpendicular to the drawing field) appears and consequently a transverse photo-emf. For example, let the drawing field be directed along the 0x axis (in a cubic crystal we choose the direction [100] as the 0x axis), and the electric vector **F** of the variable field lies in the xy plane and makes an angle φ with the direction of the drawing field. According to Eq. (1) the photocurrent in the direction 0y is given by

$$j_y^{\mathrm{ph}} = 2\gamma_{44} E F_x F_y = \gamma_{44} E F^2 \sin 2\varphi. \qquad (3)$$

This photocurrent creates a transverse emf

$$\mathscr{E}_y^{\mathrm{ph}} = -(\gamma_{44}/\sigma) l_y E F^2 \sin 2\varphi. \qquad (4)$$

Here σ is the dark electrical conductivity, and l_y is the thickness of the sample in the 0y direction.

It is easy to see that a transverse photo-emf arises only when the radiation is polarized. Its absolute value is maximal when $\varphi = \pi/4$ or $3\pi/4$, and is equal to zero when the electric vector of the variable field is either parallel or perpendicular to the drawing field.

The dependence of the photoconductivity, i.e., the photocurrent in the direction of the drawing field, on the polarization of the radiation is another anisotropic photoelectric effect. This effect is discussed in articles[2,3]. According to Eq. (1) the photoconductivity is given by

$$j_x^{\mathrm{ph}} = \{(\gamma_{11} + \gamma_{12})/2 + [(\gamma_{11} - \gamma_{12})/2] \cos 2\varphi\} E F^2. \qquad (5)$$

In an isotropic medium, as follows from relation (2), the formation of a transverse photo-emf is directly related to dependence of the photoconductivity on the angle between the direction of polarization of the radiation and the constant field. Below we shall consider specific mechanisms for the formation of anisotropic photoelectric effects.

2. PHOTOELECTRIC EFFECTS IN MANY-VALLEY SEMICONDUCTORS

In many-valley semiconductors upon exposure to polarized light which ionizes the impurity atoms, the probabilities for electron transitions into the various valleys in general turn out to be different. Therefore, the populations of the valleys are not identical. Due to the anisotropy of the mobility of the current carriers in each valley, the electrical conductivity of such a semiconductor is anisotropic. Therefore the drawing field causes the appearance of a transverse photocurrent (or transverse photo-emf), and the photoconductivity depends on the polarization of the radiation.

Let us consider a hydrogen-like donor atom in a many-valley semiconductor. Taking the corrections to the effective mass approximation into account, the ground state of an electron on a donor is a singlet state.[4,5] In particular, for an As atom in Ge the energy of excitation from the ground state to the nearest level (in Ge this is a triplet state) amounts to 4.2×10^{-3} eV.[5] For simplicity we shall assume that all neutral impurity atoms are found in the ground state. Its wave function is a symmetric superposition of wave functions pertaining to the different minima of the conduction band:

$$\Psi_0(\mathbf{r}) = \frac{1}{\sqrt{\nu}} \sum_{\alpha}^{\nu} \Phi_0^{(\alpha)}(\mathbf{r}) \psi_{\mathbf{k}_\alpha}(\mathbf{r}), \qquad (6)$$

where ψ_{k_α} is the electron Bloch function, corresponding to the α^{th} minimum of the conduction band with quasi-wave vector k_α, $\Phi_0^{(\alpha)}$ is the solution of the effective mass equation for the α^{th} minimum, ν is the number of equivalent minima. Under the action of photons whose energies exceed the ionization energy of an impurity, the electrons undergo transitions to states of the continuum which are described by Coulomb functions, $\Phi_p^{(\alpha)}$.

The cross section $s(\omega)$ for photoionization of an impurity center is put together out of the partial cross sections $s^{(\alpha)}(\omega)$ corresponding to the transition of an electron to the α-th valley:

$$s(\omega)=\frac{1}{\nu}\sum_\alpha s^{(\alpha)}(\omega). \tag{7}$$

It is easy to show that

$$s^{(\alpha)}(\omega)=\sum_{i,j}(F_i/F)s_{ij}^{(\alpha)}(\omega)(F_j/F), \tag{8}$$

where

$$s_{ij}^{(\alpha)}(\omega)=\frac{4\pi^2e^2\omega}{\varkappa^{1/2}c}\sum_p\langle\Phi_0^{(\alpha)}|x_i|\Phi_p^{(\alpha)}\rangle\langle\Phi_p^{(\alpha)}|x_j|\Phi_0^{(\alpha)}\rangle\,\delta(\varepsilon_p-\varepsilon_0-\hbar\omega), \tag{9}$$

x is the coordinate operator, ϵ_0 and ϵ_p denote, respectively, the energies of the initial and final states, κ is the dielectric constant. One can diagonalize the Hermitian tensor $s_{ij}^{(\alpha)}$. In n-Ge and n-Si the equal-energy surfaces near the minima are ellipsoids of revolution. From the symmetry of the problem it follows that $s_{ij}^{(\alpha)}$ has a diagonal form in the system of principal axes for the ellipsoid of the α-th valley. We denote the components of $s_{ij}^{(\alpha)}$ in the principal-axes system by $s_\parallel$ and $s_\perp$.

Let us calculate the number of electrons in each valley for the case of steady-state illumination. In the first place it is necessary to take account of the intervalley transitions which tend to equalize the number of electrons in the valleys. Let us represent the rate of change of the electron concentration n_α in the α-th valley due to such transitions in the form

$$\sum_\beta(n_\beta-n_\alpha)/\nu\tau_{iv}, \tag{10}$$

where τ_{iv} is a relaxation time for intervalley transitions. In addition, the electrons are captured by ionized donors at a rate n_α/τ, where τ is the lifetime of the current carriers. Under steady-state conditions

$$(s^\alpha N_d^0 I)/\nu+\sum_\beta(n_\beta-n_\alpha)/\nu\tau_{iv}-n_\alpha/\tau=0, \tag{11}$$

where I is the current density of photons, N_d^0 is the concentration of neutral donors. From Eqs. (11) it follows that

$$n_\alpha=\frac{n}{\nu}\left(1+\frac{\tau_{iv}}{\tau}\right)^{-1}\left(1+\frac{\tau_{iv}}{\tau}\frac{s^{(\alpha)}}{s}\right), \tag{12}$$

where

$$n=\sum_\alpha n_\alpha=s\tau N_d^0 I.$$

The mobility of the electrons in each valley is described by a tensor whose components in the principal-axes system of the corresponding ellipsoid we denote by $\mu_\parallel$ and $\mu_\perp$. From Eq. (12) it follows that the electrical conductivity tensor for an illuminated semiconductor is given by

$$\sigma_{ij}=en\mu\left\{\delta_{ij}+\left(1+\frac{\tau}{\tau_{мд}}\right)^{-1}\left[\frac{1}{\nu\mu s}\sum_\alpha\mu_{ij}^{(\alpha)}s^{(\alpha)}-\delta_{ij}\right]\right\}. \tag{13}$$

Here μ is the average mobility of the electrons:

$$\mu=(2\mu_\perp+\mu_\parallel)/3.$$

From (8) and (13) it follows that the photoconductivity tensor is given by

$$\sigma_{ijkl}^{ph}=e\mu\tau sN_d^0(c\varkappa^{1/2}/8\pi\hbar\omega)\left\{\delta_{ij}\delta_{kl}+(1+\tau/\tau_{iv})^{-1}\cdot\frac{1}{\nu\mu s}\left[\sum_\alpha\mu_{ij}^{(\alpha)}s_{kl}^{(\alpha)}-\delta_{ij}\delta_{kl}\right]\right\}. \tag{14}$$

In n-Ge the photoconductivity tensor in a system of cubic axes has the form

$$\gamma_{11}=\gamma_{12}=e\mu\tau sN_d^0c\varkappa^{1/2}/8\pi\hbar\omega\equiv\gamma, \tag{15}$$

$$\gamma_{44}=\gamma(1+\tau/\tau_{iv})^{-1}\frac{\Delta\mu}{\mu}\frac{\Delta s}{s}, \tag{15a}$$

where

$$\Delta\mu=(\mu_\parallel-\mu_\perp)/3;\ \Delta s=(s_\parallel-s_\perp)/3. \tag{16}$$

In n-Si in a system of cubic axes we obtain

$$\gamma_{11}=\gamma\left\{1+2\left(1+\frac{\tau}{\tau_{iv}}\right)^{-1}\frac{\Delta\mu}{\mu}\frac{\Delta s}{s}\right\},$$

$$\gamma_{12}=\gamma\left\{1-\left(1+\frac{\tau}{\tau_{iv}}\right)^{-1}\frac{\Delta\mu}{\mu}\frac{\Delta s}{s}\right\},$$

$$\gamma_{44}=0. \tag{17}$$

From (15) it follows that in n-Ge when the drawing field is directed along the [100] axis, the photoconductivity does not depend on the polarization of the radiation. However, the transverse photo-emf, determined by γ_{44}, does not vanish. On the other hand, in n-Si the transverse photo-emf associated with the same direction of the field is not present, but the photoconductivity depends on the polarization of the light. Of course, if the field is not directed along the [100] axis, then in general both effects are different from zero.

If in n-Ge the static field is directed along [100], then the ratio of the transverse photo-emf ϵ to the change δV in the voltage on the current contacts of the sample, i.e., to the usual photoconductivity, is given by

$$\mathcal{E}/\delta V=-\frac{\gamma_{44}}{\nu}\frac{l_y}{l_x}\sin 2\varphi=-\left(1+\frac{\tau}{\tau_{iv}}\right)^{-1}\frac{\Delta\mu}{\mu}\frac{\Delta s}{s}\frac{l_y}{l_x}\sin 2\varphi. \tag{18}$$

In n-Ge at liquid helium temperatures, the rate of intervalley transitions is determined by impurity scattering:[6] The time τ_{iv} associated with the emission or absorption of phonons is very large. According to[6]

$$^3/_4\tau_{iv}^{-1}=A(T)(N_d^+ + N_a^-)+B(T)N_d^0. \tag{19}$$

Here N_d^+ and N_a^- denote the concentrations of ionized donors and acceptors, which are equal at low temperatures. Expression (19) differs from that given in[6] by taking into consideration the scattering of electrons by compensated acceptor impurities. We have assumed that the probability of an intervalley transition associated with scattering by an ionized acceptor is the same

as for scattering by an ionized donor. The factor 3/4 is related to the different definitions of τ_{iv} used in[6] and in formula (10).

The lifetime of the electrons is given by $\tau = [\alpha(T)N_d^+]^{-1}$, where $\alpha(T)$ is the capture coefficient. Therefore

$$\tau_{iv}/\tau = (3\alpha/4)[2A(T) - B(T) + B(T)(N_d^0 + N_d^+)/N_d^+]^{-1}. \quad (20)$$

In n-Ge at 4.2° K, $\alpha \sim 2 \times 10^{-6}$ cm^3 sec^{-1}.[7] In Ge doped with As at 20° K (no data exist for lower temperatures) A = 2×10^{-5} cm^3/sec, and B = 0.7×10^{-5} cm^3/sec.[6]

One would expect that because of the large value of the ratio m_l/m_t in n-Ge, the ratios $\Delta\mu/\mu$ and $\Delta s/s$ should be of the order of unity. With the aid of (18), (20), and the estimates given above for α, A, and B we find that the value of the transverse photo-emf in n-Ge at liquid helium temperatures amounts to approximately 1/10 to 1/100 of the ordinary photo-response (of δV).

Anisotropic photoelectric effects also appear in many-valley semiconductors in connection with self-absorption.

3. TRANSVERSE PHOTO-EMF IN AN ISOTROPIC PLASMA

We shall determine the value of γ_{44} for an isotropic plasma semiconductor (and also for a weakly-ionized, gaseous plasma) having a fixed concentration of electrons. In this case the ordinary photoconductivity is associated with "heating" of the electrons by radiation. In order to calculate the photoconductivity tensor for a plasma, it is necessary to determine the electron distribution function to third order in the field, $\mathbf{E} + \mathbf{F}\cos\omega t$. The function $n(t, \mathbf{p})$ satisfies the kinetic equation

$$\frac{\partial n(t,\mathbf{p})}{\partial t} + e(\mathbf{E} + \mathbf{F}\cos\omega t)\frac{\partial n(t,\mathbf{p})}{\partial \mathbf{p}} - S\{n(t,\mathbf{p})\} = 0, \quad (21)$$

where $\mathbf{p}$ denotes the electron momentum, $S\{n\}$ is the collision integral (see, for example,[8]). We shall confine our attention to the case when the inelasticity of the electron scattering by phonons (or by molecules in a gas) is small. Solving the kinetic equation (21) by the method of successive approximations, we find that in the first approximation with respect to the field the distribution function is given by

$$n^{(1)}(t,\mathbf{p}) = -e\frac{dn^{(0)}}{d\varepsilon}\tau_1(\varepsilon)\mathbf{v}_p \cdot \left\{\mathbf{E} + \frac{1}{2}[(1 - i\omega\tau_1)^{-1}e^{-i\omega t} + (1 + i\omega\tau_1)^{-1}e^{i\omega t}]\mathbf{F}\right\}. \quad (22)$$

Here τ_l is the relaxation time for the distribution-function spherical harmonic $n_{lm}(\epsilon)$, $\mathbf{v}_p$ is the electron velocity, $n^{(0)}$ is the equilibrium distribution function.

From the equation for the second-order approximation to the distribution function,

$$\frac{\partial n^{(2)}}{\partial t} - S\{n^{(2)}\} = -e(\mathbf{E} + \mathbf{F}\cos\omega t)\frac{\partial n^{(1)}}{\partial \mathbf{p}} \quad (23)$$

it follows that $n^{(2)}$ contains harmonics with $l = 0$ and $l = 2$. The zero harmonic describes the "heating" of an electron gas by radiation, the second harmonic describes a "quadrupole" deformation of the distribution function. The equations for $n_0^{(2)}$ and $n_{2m}^{(2)}$ follow from Eq. (23):

$$\frac{\partial n_0^{(2)}}{\partial t} - \frac{1}{N(\varepsilon)}\frac{d}{d\varepsilon}\left\{\frac{\varepsilon N(\varepsilon)}{\tau_0(\varepsilon)}\left(T\frac{dn_0^{(2)}}{d\varepsilon} + n_0^{(2)}\right)\right\} = \frac{4}{3}\frac{e^2}{m}\mathbf{EF}\cos\omega t\frac{1}{N(\varepsilon)}\frac{d}{d\varepsilon}\left\{N(\varepsilon)\varepsilon\tau_1\frac{dn^{(0)}}{d\varepsilon}\right\}, \quad (24)$$

$$\frac{\partial n_2^{(2)}}{\partial t} + \frac{n_2^{(2)}}{\tau_2} = \left\{e^2\,\mathrm{Re}\frac{d}{d\varepsilon}\left[\frac{dn^{(0)}}{d\varepsilon}\tau_1(1 + (1 + i\omega\tau_1)^{-1})\right]e^{i\omega t}E_iF_k + \frac{e^2}{2}\frac{d}{d\varepsilon}\left[\frac{dn^{(0)}}{d\varepsilon}\tau_1(1 + \omega^2\tau_1^2)^{-1}\right]F_iF_k\right\}\left(v_iv_k - \frac{v^2}{3}\delta_{ik}\right). \quad (25)$$

Here $N(\epsilon)$ denotes the density of states. Only terms which at the end of the calculation give a contribution to the time-independent part of $n^{(3)}$ are left in the right-hand side of Eq. (25). In order to determine the photocurrent we only need the constant component of $n^{(3)}$ having the symmetry of the spherical harmonic function with $l = 1$:

$$n_{1m}^{(3)} = \tau_1\overline{\left\{-\frac{\partial n^{(2)}}{\partial \mathbf{p}}e(\mathbf{E} + \mathbf{F}\cos\omega t)\right\}}_{l=1}. \quad (26)$$

The equation for $n_0^{(2)}$ can be integrated only for $\omega\tau_0 \ll 1$ or $\omega\tau_0 \gg 1$. For $\omega\tau_0 \ll 1$ one finds

$$n_0^{(2)} = \frac{2e^2n^{(0)}}{3mT^2}(\mathbf{E} + \mathbf{F}\cos\omega t)^2\left\{\int_0^\varepsilon d\varepsilon\tau_0(\varepsilon)\tau_1(\varepsilon) - \int_0^\infty d\varepsilon f(\varepsilon)\tau_0(\varepsilon)\tau_1(\varepsilon)\right\}, \quad (27)$$

where

$$f(\varepsilon) = \frac{1}{n}\int_\varepsilon^\infty d\varepsilon N(\varepsilon)n^{(0)}(\varepsilon).$$

For $\omega\tau_0 \gg 1$ that part of $n_0^{(2)}(\omega)$ which oscillates in phase with the field $\mathbf{F}(t)$ (it gives a contribution to γ_{44}) amounts to a quantity of the order of $n_0^{(2)}(0)/\omega^2\tau_0^2$.

It is easy to integrate Eq. (25) for arbitrary ω. Substituting the result of the integration into Eq. (26), and then substituting the resulting expression into the formula for the current density, we find

$$\gamma_{44}(\omega) = \gamma_{44}^{(0)}(\omega) + \gamma_{44}^{(2)}(\omega), \quad (28)$$

where $\gamma_{44}^{(0)}(\omega)$ and $\gamma_{44}^{(2)}(\omega)$ correspond to the functions $n_0^{(2)}$ and $n_{2m}^{(2)}$:

$$\gamma_{44}^{(0)} = \frac{4e^4}{9m^2T^2}\int\frac{d^3p}{(2\pi)^3}\varepsilon\tau_1\frac{d}{d\varepsilon}\left\{n^{(0)}\left[\int_0^\infty d\varepsilon f(\varepsilon)\tau_0\tau_1 - \int_0^\varepsilon d\varepsilon\tau_0\tau_1\right]\right\}, \quad \omega\tau_0 \ll 1,$$

$$\gamma_{44}^{(0)}(\omega) \sim \gamma_{44}^{(0)}(0)/\omega^2\tau_0^2, \quad \omega\tau_0 \gg 1,$$

$$\gamma_{44}^{(2)} = \frac{2e^4}{3m^2}\int\frac{d^3p}{(2\pi)^3}\tau_1\varepsilon\left\{\left(\frac{4}{5}\frac{dM}{d\varepsilon} + \frac{7}{30}\frac{dN}{d\varepsilon}\right)\varepsilon + M + \frac{1}{6}N\right\},$$

$$M = \tau_2 d/d\varepsilon\left(\tau_1\frac{dn^{(0)}/d\varepsilon}{1 + \omega^2\tau_1^2}\right),$$

$$N = \mathrm{Re}\,\tau_2(1 + i\omega\tau_2)^{-1}\frac{d}{d\varepsilon}\left\{\frac{dn^{(0)}}{d\varepsilon}\tau_1[1 + (1 + i\omega\tau_1)^{-1}]\right\}. \quad (29)$$

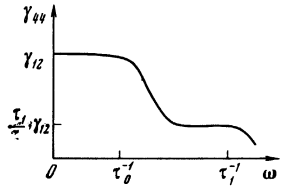
One can call the contribution to γ_{44} coming from $n_0^{(2)}$, i.e., $\gamma_{44}^{(0)}$, the "heating" contribution. It is related to the fact that the energy distribution function $n_0^{(2)}$ oscillates at the frequency of the radiation and has an amplitude proportional to $(\mathbf{E} \cdot \mathbf{F})$. This part of $\gamma_{44}(\omega)$ starts to fall off with frequency at $\omega \sim \tau_0^{-1}$.

A calculation similar to the one carried out above shows that for $\omega\tau_1 \ll 1$

$$\gamma_{12}(\omega) - \gamma_{12}^{(0)}(0) \sim e\mu\tau_0\sigma/T,$$

where $\mu = e\tau_1/m$ denotes the mobility of the current carriers.

The contribution to γ_{44} coming from $n_{2m}^{(2)}$, i.e., $\gamma_{44}^{(2)}$, which is related to a deformation of the momentum distribution function, can be referred to as the "deformation" contribution. For $\omega\tau_1 \ll 1$ we have $\gamma_{44}^{(2)} \sim e\mu\tau_1\sigma/T$, that is, this quantity is smaller than $\gamma_{44}^{(0)}(0)$ by the factor τ_1/τ_0. Therefore at low frequencies ($\omega\tau_0 \ll 1$) the value of γ_{44} is determined by the "heating" contribution, but for $\omega \gtrsim (\tau_0\tau_1)^{-1/2}$ the deformation part is dominant. The frequency dependence of $\gamma_{44}(\omega)$ in an isotropic plasma is schematically shown in the accompanying Figure.



Frequency dependence of γ_{44} for an isotropic electron gas (schematical).

Thus, in the low frequency region ($\omega \lesssim \tau_0^{-1}$) the effects determined by the components of the photoconductivity tensor (the transverse photo-emf, the dependence of the photoconductivity on polarization) are not small; they are of the same order as the ordinary "heating" photoconductivity. At higher frequencies ($\tau_0^{-1} \ll \omega < \tau_1^{-1}$) the anisotropic photoelectric effects are smaller than the ordinary photoconductivity by the factor τ_1/τ_0.

The electron gas in n-InSb is an example of an isotropic plasma. At liquid helium temperatures $\tau_0 \sim 5 \times 10^{-7}$ sec and $\tau_1 \sim 5 \times 10^{-13}$ sec. Therefore over the entire uhf band $\omega\tau_0 \gg 1$, and anisotropic photoelectric effects are negligible ($\sim 10^{-6}$ times the plasma photoconductivity). At liquid nitrogen temperatures optical phonons play an appreciable role in scattering the momentum of the electrons in n-InSb. Due to the large inelasticity of the scattering by optical phonons, one would expect the times for the scattering of energy and momentum to be of the same order ($\sim 10^{-11}$ sec). Therefore, up to frequencies $\sim 10^{11}$ Hz the anisotropic photoelectric effects should be of the same order as the ordinary plasma photoconductivity of n-InSb.

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