05.2

Anisotropic photoconductivity excited in a semiconductor by two-frequency optical radiation

© P.A. Ziaziulia[1], V.L. Malevich[2,3]

[1] Belarusian State University, Minsk, Belarus
[2] B.I. Stepanov Institute of Physics, National Academy of Sciences of Belarus, Minsk, Belarus
[3] Belarusian State University of Informatics and Radioelectronics, Minsk, Belarus
E-mail: zezyulya@bsu.by

Received July 19, 2022
Revised September 1, 2022
Accepted September 14, 2022.

Anisotropic photoconductivity at a difference frequency, excited in a semiconductor by linearly polarized two-frequency optical radiation, is considered. The anisotropy of the photoconductivity arises due to the optical alignment of photoexcited electrons momenta and dependence of their effective mass and momentum relaxation time on energy. It is shown that the contribution of the anisotropic photoconductivity to the photocurrent at the difference frequency lying in the terahertz frequency range can be comparable with that of the isotropic photoconductivity. This effect can manifest itself in photoconductive antennae, devices used to generate terahertz radiation.

Keywords: Anisotropic photoconductivity, photomixing, terahertz radiation, photoconductive antenna.

DOI: 10.21883/TPL.2022.11.54881.19313

One of the methods for generating continuous terahertz (THz) radiation is the use of the photomixing effect under interband optical excitation of a semiconductor by two light beams, the frequency difference of which lies in the THz region [1,2]. The photomixing effect arises due to the nonlinear nature of photoconductivity and in the presence of a constant bias leads to the generation of a photocurrent at the beat frequency. THz-emitter based on the photomixing effect, called a photoconductive antenna, is a photosensitive semiconductor layer, on the surface of which a system of metal electrodes is formed, which plays the role of a radiating antenna, and also serves to supply a bias voltage. For the effective operation of the photoconductive antenna as a THz emitter, it is necessary that the lifetime of nonequilibrium charge carriers, which determines the photoresponse time of the semiconductor, be less than 1 ps. Therefore, in photoantennas used to generate THz radiation, it is necessary to use specially grown semiconductor layers with subpicosecond lifetimes of nonequilibrium current carriers [1,2].

In this paper we review the anisotropic photoconductivity in the semiconductor excited by two-frequency optical radiation, as well as its contribution to the photocurrent at the beat frequency corresponding to the THz region. The anisotropic photoconductivity for interband optical transitions is due to the anisotropy of the momenta distribution of electrons excited by polarized light and the dependence of their momentum relaxation time and effective mass on energy [3–5]. The response time of anisotropic photoconductivity is determined by the electron momentum relaxation time, which is usually 200−300 fs.

Consider a cubic semiconductor excited by two-frequency linearly polarized optical radiation $\mathbf{E}(t) = \mathbf{E}_1 \cos\omega_+ t + \mathbf{E}_2 \cos\omega_- t$ with close frequencies $\omega_\pm = \omega \pm \Omega/2 > \varepsilon_g/\hbar$ (ε_g is semiconductor band gap). A constant electric field $\mathbf{F}$ applied to the semiconductor is directed along the z axis, the vectors $\mathbf{E}_1$ and $\mathbf{E}_2$ are parallel to each other and lie in the xz plane at an angle γ to the z axis. The concentration of photoexcited electrons and holes is proportional to the intensity of the incident optical radiation and includes two terms, one of which is time independent, and the second varies at the difference frequency Ω, which is assumed lying in the THz-range. Here we are interested in the photocurrent component at the beat frequency, since this component excites THz radiation in the photoconductive antenna.

The kinetic equation for the component of the distribution function $f_\mathbf{p}$ of photoexcited electrons over momenta $\mathbf{p}$ at the difference frequency in the constant electric field has the form

$$i\Omega f_\mathbf{p} - e\mathbf{F}\frac{\partial f_\mathbf{p}}{\partial \mathbf{p}} = G_\mathbf{p}(t) - \frac{f_0(\varepsilon_\mathbf{p})}{\tau_0} + I_c(f_\mathbf{p}), \qquad (1)$$

$$G_\mathbf{p}(t) = \frac{\alpha I \delta(\varepsilon_0 - \varepsilon_\mathbf{p})}{\hbar\omega g(\varepsilon_\mathbf{p})}\big(1 - P_2(\cos\beta)\big), \qquad (2)$$

where τ_0 is recombination time of non-equilibrium electrons, $I_c(f_\mathbf{p})$ is collision integral, $f_0(\varepsilon_\mathbf{p})$ is momentum-symmetric part of the photoelectron distribution function, $P_2(x)$ is second-order Legendre polynomial, β is angle between direction of electron momentum and polarization vector of optical radiation, $\delta(x)$ is delta function, I is intensity amplitude of exciting radiation at frequency Ω, $g(\varepsilon_\mathbf{p})$ is electron states density in the conduction band, $\mathbf{v}_\mathbf{p} = \partial\varepsilon_\mathbf{p}/\partial\mathbf{p}$ is speed

of electron with energy $\varepsilon_{\mathbf{p}}$, α is coefficient of interband light absorption. Expression (2) describes the generation of non-equilibrium electrons during interband transitions from the heavy hole subband of the valence band (the contribution of transitions from the light hole subband is neglected) [3]. The energy of the photoexcited electron ε_0 is determined from the relation $\varepsilon_{\mathbf{p}} + \varepsilon_{v,\mathbf{p}} = \hbar\omega - \varepsilon_g$ ($\varepsilon_{v,\mathbf{p}}$ is dispersion law for heavy holes).

Kinetic equation (1) will be solved in the linear approximation over the field $\mathbf{F}$ and in the quadratic approximation over the light wave field, using the standard procedure [6] of the electron distribution function expansion by spherical functions Y_{lm}

$$f_{\mathbf{p}} = f_0(\varepsilon_{\mathbf{p}}) + \sum_{l=1,2} \sum_{m=-l}^{+l} f_{l,m}(\varepsilon_{\mathbf{p}}) Y_{lm}(\theta, \varphi), \qquad (3)$$

where θ and φ are the polar and azimuth angles of the vector $\mathbf{p}$ in the spherical coordinate system. For the case of quasi-elastic scattering considered here, the collision integral can be represented as

$$I_c(f_{\mathbf{p}}) = -\frac{1}{\tau_1(\varepsilon_{\mathbf{p}})} \sum_{m=-1}^{+1} f_{1,m}(\varepsilon_{\mathbf{p}}) Y_{1m}(\theta, \varphi) - \frac{1}{\tau_2(\varepsilon_{\mathbf{p}})} \sum_{m=-2}^{+2} f_{2,m}(\varepsilon_{\mathbf{p}}) Y_{2m}(\theta, \varphi), \qquad (4)$$

where τ_1 and τ_2 are the relaxation times of the first and second spherical harmonics of the distribution function. Using the addition theorem for spherical functions, we write

$$P_2(\cos\beta) = \frac{4\pi}{5} \sum_{m=-2}^{+2} Y_{2m}(\theta, \varphi) Y^*_{2m}(\gamma, 0). \qquad (5)$$

Substituting expressions (2)–(5) into (1), we obtain the equation for the coefficients of the spherical function expansion of the distribution function. Multiplying it by $Y^*_{lm}(\theta, \varphi)$ and integrating over the angles, we find a chain of equations for the coefficients $f_{l,m}$. In a linear approximation over constant electric field, from this system of equations we obtain expressions for the expansion coefficients that determine the photocurrent:

$$f_{1,0} = 2\sqrt{\frac{\pi}{3}} \frac{\alpha I}{\hbar\omega} \frac{eF\tau_1 v_{\mathbf{p}}}{1+i\Omega\tau_1} \left[\frac{\tau_0}{1+i\Omega\tau_0} \frac{d}{d\varepsilon_{\mathbf{p}}} \left(\frac{\delta(\varepsilon_0 - \varepsilon_{\mathbf{p}})}{g(\varepsilon_{\mathbf{p}})} \right) - \frac{2}{5} P_2(\cos\gamma) U(\varepsilon_{\mathbf{p}}) \right], \qquad (6)$$

$$f_{1,\pm1} = \pm \frac{\sqrt{6\pi}}{10} \frac{\alpha I \sin 2\gamma}{\hbar\omega} \frac{eF\tau_1 v_{\mathbf{p}}}{1+i\Omega\tau_1} U(\varepsilon_{\mathbf{p}}), \qquad (7)$$

$$U(\varepsilon_p) = \frac{d}{d\varepsilon_{\mathbf{p}}} \left[\frac{\tau_2}{1+i\Omega\tau_2} \frac{\delta(\varepsilon_0 - \varepsilon_{\mathbf{p}})}{g(\varepsilon_{\mathbf{p}})} \right] + 3\frac{\tau_2}{1+i\Omega\tau_2} \frac{\delta(\varepsilon_0 - \varepsilon_{\mathbf{p}})}{p g(\varepsilon_{\mathbf{p}}) v_{\mathbf{p}}}. \qquad (8)$$

After substituting (6) and (7) into the expression for the current, and integrating over momenta for the longitudinal j_z and transverse j_x components of the photocurrent, we obtain

$$j_z = \frac{e^2 F}{3\pi^2} \frac{\alpha I}{\hbar^4 \omega g(\varepsilon_0)} \left\{ \frac{\tau_0}{1+i\Omega\tau_0} \frac{d}{d\varepsilon_{\mathbf{p}}} \left(p^2 v_{\mathbf{p}} \frac{\tau_1}{1+i\Omega\tau_1} \right) - \frac{2}{5} P_2(\cos\gamma) \frac{\tau_2}{1+i\Omega\tau_2} V(\varepsilon_{\mathbf{p}}) \right\} \Bigg|_{\varepsilon_{\mathbf{p}}=\varepsilon_0}, \qquad (9)$$

$$j_x = -\frac{e^2 F}{10\pi^2} \frac{\alpha I \sin 2\gamma}{\hbar^4 \omega g(\varepsilon_0)} \frac{\tau_2}{1+i\Omega\tau_2} V(\varepsilon_{\mathbf{p}}) \Bigg|_{\varepsilon_{\mathbf{p}}=\varepsilon_0}, \qquad (10)$$

$$V(\varepsilon_{\mathbf{p}}) = p^3 \frac{d}{d\varepsilon_{\mathbf{p}}} \left(\frac{v_{\mathbf{p}}}{p} \frac{\tau_1}{1+i\Omega\tau_1} \right) = \frac{\tau_1}{1+i\Omega\tau_1} p^3 \frac{d}{d\varepsilon_{\mathbf{p}}} \left(\frac{v_{\mathbf{p}}}{p} \right) + p^2 v_{\mathbf{p}} \frac{d}{d\varepsilon_{\mathbf{p}}} \left(\frac{\tau_1}{1+i\Omega\tau_1} \right). \qquad (11)$$

The first term in expression (9), containing the lifetime of non-equilibrium electrons τ_0, describes isotropic photoconductivity. The anisotropy of the photoexcited electrons distribution over momenta leads to the appearance of the second term in (9), which depends on the angle γ, as well as of the transverse component of the photocurrent (10) perpendicular to the pulling field $\mathbf{F}$. Note that expressions (9)–(11) were obtained neglecting energy relaxation and, therefore, are valid under the condition $\Omega\tau_\varepsilon \gg 1$ (τ_ε is energy relaxation time of electrons).

The anisotropic part of the photoconductivity is determined by the function $V(\varepsilon_{\mathbf{p}})$, which is nonzero only when the effective mass and (or) the momentum relaxation time of the electrons depend on the energy. The first term in (11) describes the contribution to the anisotropic photoconductivity due to the nonparabolicity of the electronic spectrum, and the second term due to the energy dependence of the momentum relaxation time.

The function $V(\varepsilon_{\mathbf{p}})$ is determined by the energy derivatives of the effective mass and momentum relaxation time of electrons. Therefore, in the case of a threshold dependence of these parameters on energy, for example, when the energy of the photoelectron lies near the threshold of intervalley transitions or the threshold of the optical phonon emission, one can expect the anisotropic component of photoconductivity increasing. In the paper [7] the spectral dependence of the anisotropic photoconductivity was investigated by the Monte-Carlo method, and it was shown that the anisotropic photoconductivity reaches its maximum when the energy of the photoexcited electron lies near the threshold of intervalley transitions. In typical semiconductors, the relation $\tau_0 \gg \tau_1, \tau_2$ is usually satisfied. In the low-frequency region ($\Omega\tau_0 \ll 1$) the anisotropic addition to photoconductivity at the difference frequency is small, since its relation to ordinary photoconductivity is characterized by the parameter $\tau_2(\varepsilon_0)/\tau_0 \ll 1$. However, in the THz-frequency region $\Omega\tau_0 \gg 1$, the ratio of the

anisotropic part of the photoconductivity to the isotropic part is determined by the parameter $\Omega\tau_2(\varepsilon_0)$, the value of which may be about unity. So, at the terahertz difference frequency the anisotropic component of the photoconductivity can be comparable with the usual isotropic one. This conclusion is also valid for semiconductor layers specially grown by molecular beam epitaxy and used in photoconductive antennas, since for them in the THz frequency range $\tau_0 \sim \Omega^{-1}$.

Conflict of interest

The authors declare that they have no conflict of interest.

References

[1] E. Peytavit, G. Ducournau, J.-F. Lampin, in *Fundamentals of terahertz devices and applications*, ed. by D. Pavlidis (John Wiley & Sons, Hoboken, N.J., 2021), p. 137–186. DOI: 10.1002/9781119460749
[2] N. Burford, M. El-Shenawee, Opt. Eng., **56** (1), 010901 (2017). DOI: 10.1117/1.OE.56.1.010901
[3] V.I. Belinicher, V.N. Novikov, Sov. Phys. Semicond., **15** (10), 1138 (1981).
[4] M.I. Karaman, V.P. Mushinskii, G.M. Shmelev, Sov. Phys. Tech. Phys., **28**, 730 (1983).
[5] S.Kh. Esayan, E.L. Ivchenko, V.V. Lemanov, A.Yu. Maksimov, JETP Lett., **40** (11), 1290 (1984).
[6] I.M. Dykman, P.M. Tomchuk, *Yavleniya perenosa i fluktuatsii v poluprovodnikakh* (Nauk. dumka, Kiev, 1981). (in Russian)
[7] Y.V. Malevich, R. Adomavičius, A. Krotkus, V.L. Malevich, J. Appl. Phys., **115** (7), 073103 (2014). DOI: 10.1063/1.4865961