Semiconductors, 2024, Vol. 58, No. 9

Anisotropy of polarization of interband photoluminescence in n-InAs induced by electric field

© R.B. Adamov, M.Ya. Vinnichenko¶, N.Yu. Kharin, D.A. Karaulov, D.A. Firsov

Peter the Great St. Petersburg Polytechnic University,
195251 St. Petersburg, Russia

¶ E-mail: mvin@spbstu.ru

Received September 30, 2024
Revised November 7, 2024
Accepted November 7, 2024

The degree of linear polarization of interband photoluminescence in an InAs crystal doped with donors in an electric field was calculated. The polarization anisotropy arises due to the anisotropy of the electron distribution function over states in momentum space, associated with the electron drift in the electric field, and the dependence of the optical matrix elements on the angle between the polarization vector and the electron wave vector. A quasi-equilibrium distribution function shifted in velocity space was used. The electron temperature was determined from the power balance equation. The effect of nonequilibrium phonon accumulation was taken into account in calculating the rate of energy loss by hot electrons. The nonparabolicity of the conduction band was taken into account using the Kane dispersion law.

Keywords: distribution function anisotropy, interband photoluminescence, polarization anisotropy, electron temperature.

DOI: 10.61011/SC.2024.09.59920.7123

1. Introduction

Photoluminescence (PL) is a powerful non-destructive tool for the characterization of semiconductors and semiconductor nanostructures [1]. The analysis of photoluminescence spectra makes it possible to determine such important characteristics as the energy spectrum of charge carriers [2], the binding energy of excitons [3], and the lifetimes of nonequilibrium charge carriers [4]. The characteristic times of the electron-electron and electron-phonon interactions in graphene were determined using the analysis of photoluminescence spectra [5]. The temperature maps of the crystal lattice and hot electrons of were simultaneously determined in GaN HEMT devices [6]. The study of photoluminescence during interband optical pumping is the initial stage of the development of many optoelectronic devices (quantum cascade and injection lasers, LEDs, photodetectors, single photon sources, etc.).

The study of photoluminescence and other optical phenomena under highly nonequilibrium conditions is of particular interest. The study of optical phenomena related to nonequilibrium charge carriers is of considerable interest, since many optoelectronic devices (cascade and injection semiconductor lasers, cascade photodetectors, radiation modulators) operate under conditions of heating of charge carriers. Hot electron detectors of radiation based on metal−semiconductor plasmonic structures were created [7–9]. The sensitivity of such detectors is still low compared to classical devices based on InAs, GaAs and CdHgTe, however, there are prospects for its increase [7].

The absorption of radiation with a photon energy significantly exceeding the band gap is one of the ways to create nonequilibrium distributions of charge carriers. The nonequilibrium distribution of photoexcited electrons by states, which is formed in the processes of thermalization and recombination, is theoretically considered in Ref. [10], taking into account the heating of electrons, holes, and phonons.

The so-called „hot photoluminescence" [11] is a striking example of the anisotropic effect when electrons are heated by light. It can be observed if the photon energy of linearly polarized exciting radiation significantly exceeds the band gap of a semiconductor. In this case, the momenta of nonequilibrium electrons are aligned at the initial stage of energy relaxation, due to the dependence of the optical matrix element on the angle between the electron momentum and the radiation polarization vector. The anisotropic distribution of nonequilibrium electrons leads to anisotropy of the polarization of recombination radiation. Such anisotropy is observed for the short-wavelength edge of the radiation spectrum.

An anisotropic momentum distribution of nonequilibrium electrons is created by anisotropic excitation in the case of hot photoluminescence. The momentum distribution of electrons becomes symmetrical during the time of energy relaxation to the bottom of the conduction band, and the linear polarization of photoluminescence disappears. An anisotropic momentum distribution of thermalized electrons can be formed by applying an electric field that causes heating and electron drift. The drift of electrons in an electric field means that they have an average drift momentum directed along the electric field. The intensity of recombination radiation polarized parallel and perpendicular to the applied field will vary in accordance with the dependence of the optical matrix elements on the angle

between the electron momentum and the polarization vector.

The anisotropy of the polarization of luminescence of hot electrons in an electric field should be most noticeable in semiconductors with high mobility. Previously, this effect was observed in *n*-InSb [12], but this phenomenon and its spectral dependence have not been studied in detail. Photoluminescence in *n*-InN under conditions of an applied electric field was studied in Ref. [13]. However, only the effect of heating of electrons in an electric field on the luminescence spectrum was analyzed in this paper, and the polarization characteristics of PL were not studied. The results of an experimental study of the infrared photoluminescence spectra of a doped epitaxial layer of gallium arsenide under conditions of heating of charge carriers by an electric field are presented in Ref. [14]. The anisotropy of the polarization of photoluminescence radiation, which occurs due to the anisotropy of the hot electron distribution function and the angular dependence of the interband optical matrix element in the momentum space, was discovered and studied. It was shown that the experimental data are consistent with the theoretical model (calculations were performed in the parabolic approximation).

The effect of the nonequilibrium distribution of electrons over states on the optical characteristics of a semiconductor associated with interband transitions of charge carriers was also studied in Ref. [15], which demonstrated the effect of electron drift in *n*-InSb in an electric field on the interband absorption of radiation of two polarizations — along of the applied field and perpendicular to it. The dependence of radiation absorption on its polarization was described in the framework of the parabolic approximation.

This paper presents the results of a detailed calculation of the anisotropy of luminescence polarization in a high-mobility bulk semiconductor *n*-InAs. The calculation takes into account the nonparabolicity of the conduction band and the effect of accumulation of nonequilibrium optical phonons. It should be noted that the ability to control the polarization of photoluminescence radiation is of interest for the development of radiation sources with a given polarization.

2. Determination of the temperature of hot electrons

All calculations in this paper are performed for a bulk semiconductor *n*-InAs doped to the level of $n = 2.2 \cdot 10^{16}\,\mathrm{cm}^{-3}$ with the lattice temperature of $T_0 = 90\,\mathrm{K}$. The anisotropic deformation of the distribution function in the momentum space is a necessary condition for observing the anisotropy of photoluminescence polarization in a cubic semiconductor in an electric field. Electron-electron collisions effectively redistribute the energy coming from the electric field between the electrons at sufficiently high concentrations of free electrons. In this case, it is possible to use the quasi-equilibrium Fermi–Dirac or Maxwell–Boltzmann $f_e(\mathbf{k})$ distribution function ($\mathbf{k}$ — electron wave vector) with the electron temperature $T_e > T_0$ and drift velocity $\mathbf{v}_{dr}$ used as parameters [16]. The drift velocity can be determined experimentally by analyzing the current-voltage characteristic.

Let's find the dependence of the electron temperature on the applied electric field. An external electric field leads to an increase of the average energy of a free electron with a velocity of $e\mu_e\mathbf{F}^2$, where e — carrier charge, $\mathbf{F}$ — applied electric field, μ_e — the mobility of electrons. The steady state is achieved when the average rate of energy gain by electrons is balanced by the average rate of its transfer to the phonon subsystem $\langle dE/dt\rangle$:

$$e\mu_e\mathbf{F}^2 = \langle dE/dt\rangle. \tag{1}$$

The average rate of energy loss by an electron is determined by averaging the rate of energy loss dE/dt, depending on the scattering mechanism, over an ensemble of electrons. The scattering over longitudinal optical phonons is the dominant mechanism of free electron energy relaxation in InAs at liquid nitrogen temperatures and not too high electron concentrations [17]. The effect of phonon accumulation becomes noticeable due to the finite lifetime of nonequilibrium phonons emitted by electrons with sufficiently high energy, which should be taken into account in calculations. To do this, it is convenient to use a different approach to calculating the average energy loss rate based on using the rate of change in the number of phonons due to scattering.

The number of phonons is determined by their distribution function (DF) $N_{\mathbf{q}}$. The DF of optical phonons corresponds to the Bose–Einstein statistics in the equilibrium state:

$$N_{\mathbf{q}0} = \left(\exp\left(\frac{\hbar\omega_{\mathrm{LO}}}{k_{\mathrm{B}}T_0}\right) - 1\right)^{-1}, \tag{2}$$

where $\hbar\omega_{\mathrm{LO}} = 29.8\,\mathrm{meV}$ [18] — the energy of the longitudinal optical phonon in InAs, k_{B} — the Boltzmann constant.

The DF can be represented as the sum of the equilibrium DF and the nonequilibrium additive when taking into account the accumulation of LO-phonons:

$$N_{\mathbf{q}} = N_{\mathbf{q}0} + \frac{dN_{\mathbf{q}}}{dt}\,\tau_q^{\mathrm{LO}}, \tag{3}$$

where τ_q^{LO} is the lifetime of nonequilibrium LO-phonons. The lifetime of nonequilibrium optical phonons in InAs is determined by the decay of a long-wavelength optical phonon into two short-wavelength acoustic phonons [19]. We used the value $\tau_q^{\mathrm{LO}} = 5.8\,\mathrm{ps}$ for calculations, obtained in Ref. [19] for the lattice temperature of 77 K. The lifetime will slightly decrease with an increase in the lattice temperature to 90 K, which should not significantly change the results obtained.

Now it is possible to write the average rate of energy loss by an electron as the product of the rate of change of the phonon DF $dN_{\mathbf{q}}/dt$ by the phonon energy $\hbar\omega_{\mathrm{LO}}$, summed

over all phonon wave vectors $\mathbf{q}$ and divided by the number of electrons nV, where V — normalization volume:

$$\langle dE/dt\rangle = \frac{\sum_{\mathbf{q}} \hbar\omega_{\mathrm{LO}} \frac{dN_{\mathbf{q}}}{dt}}{nV}. \tag{4}$$

The rate of change of the number of phonons $\frac{dN_{\mathbf{q}}}{dt}$ can be determined using perturbation theory [16]:

$$\frac{dN_{\mathbf{q}}}{dt} = \frac{2\pi}{\hbar}\sum_{k}\big\{|\langle \mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 \times f_e(\mathbf{k}+\mathbf{q})\delta(E_e(\mathbf{k}+\mathbf{q}) - E_e(\mathbf{k}) - \hbar\omega_{\mathrm{LO}}) - |\langle \mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2 \times f_e(\mathbf{k})\delta(E_e(\mathbf{k}) - E_e(\mathbf{k}+\mathbf{q}) + \hbar\omega_{\mathrm{LO}})\big\}, \tag{5}$$

where the first term is related to the emission of phonons by free electrons, and the second term is related to their absorption; Dirac's δ-function reflects the law of conservation of energy.

When calculating $dN_{\mathbf{q}}/dt$ according to (5) in this paper, we take into account the nonparabolicity of the dispersion law for the conduction band using the Kane model [20]:

$$E_e(\mathbf{k}) = \frac{1}{2\alpha}\left(\sqrt{1+\frac{2\alpha\hbar^2k^2}{m_e}} - 1\right), \tag{6}$$

where $\alpha = \frac{1}{E_g}\left(1-\frac{m_e}{m_0}\right)^2$ — nonparabolicity parameter [21], $E_g = 0.405\,\mathrm{meV}$ — InAs band gap at 90 K [22], m_0 — mass of a free electron, $m_e = 0.026m_0$ — effective mass of electrons at the point Γ of the conduction band [22]. Using the Kane model to describe the law of dispersion of the InAs conduction band in the equilibrium case is a generally accepted approach. A relatively low concentration of free electrons is used in this paper, which occupy states near the bottom of the conduction band, so that interaction with distant bands can be neglected. The shift of the distribution function relative to the Brillouin zone is also small in electric fields not exceeding 1000 V/cm.

We use matrix elements corresponding to the parabolic approximation in (5) to simplify calculations. In this case, taking into account the screening effects, the squares of the modules of the matrix elements associated with the emission and absorption of phonons look as follows [17]:

$$|\langle \mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 = \frac{2\pi\hbar^2 eE_0}{Vm_eq^2}\,(1+(qr_{\mathrm{D}})^{-2})^{-2}(N_{\mathbf{q}}+1), \tag{7}$$

$$|\langle \mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2 = \frac{2\pi\hbar^2 eE_0}{Vm_eq^2}\,(1+(qr_{\mathrm{D}})^{-2})^{-2}N_{\mathbf{q}}, \tag{8}$$

where $E_0 = em_e\hbar\omega_{\mathrm{LO}}(\varepsilon_\infty^{-1} - \varepsilon_0^{-1})/\hbar^2$ — the interaction constant of an electron with a phonon, $\varepsilon_\infty = 11.91$ and $\varepsilon_0 = 14.55$ — high-frequency and low-frequency permittivity [18], $r_{\mathrm{D}} = (\varepsilon_\infty k_{\mathrm{B}}T_e/4\pi e^2 n)^{1/2}$ — Debye screening radius for a non-degenerate electron gas.

The electron gas is non-degenerate for the considered electron concentration n. If we do not take into account the weak anisotropy in the phonon distribution over q associated with electron drift in the field (and, accordingly, with the anisotropy of the electron momentum distribution function), then in (5) the DF of nonequilibrium electrons can be considered symmetric in momentum space and obeying Maxwell–Boltzmann statistics:

$$f_e(\mathbf{k}) = f_{e0}(\mathbf{k}) = \frac{n}{N_c}\exp\left(-\frac{E_e(k)}{k_{\mathrm{B}}T_e}\right), \tag{9}$$

where T_e is the nonequilibrium temperature of hot electrons, and $N_c = \int_{\mathrm{ZB}} \exp\left(-\frac{E_e(k)}{k_{\mathrm{B}}T_e}\right)\frac{2d\mathbf{k}}{(2\pi)^3}$ is the effective density of states in the conduction band obtained by integration within the boundaries of the Brillouin zone (ZB).

Let's note that the Dirac's δ-function is symmetric, and let's proceed in (5) from summation over all k of the first Brillouin zone to integration:

$$\frac{dN_{\mathbf{q}}}{dt} = \frac{2\pi}{\hbar}\int_{\mathrm{ZB}}\big\{|\langle \mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 \times f_e(\mathbf{k}+\mathbf{q}) - |\langle \mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2 \times f_e(\mathbf{k})\big\}\delta\big(E_e(\mathbf{k}+\mathbf{q}) - E_e(\mathbf{k}) - \hbar\omega_{\mathrm{LO}}\big)\frac{2V d\mathbf{k}}{(2\pi)^3}, \tag{10}$$

where $\frac{2Vd\mathbf{k}}{(2\cdot\pi)^3}$ is the number of quantum states in the elementary volume $d\mathbf{k} = dk_xdk_ydk_z$, taking into account the twofold spin degeneracy. Let's move on to integration using spherical coordinates ($0 \le \theta \le \pi$ — polar angle, $0 \le \varphi \le 2\pi$ — azimuthal angle, k — wave vector modulus, d — the period of the crystal lattice):

$$\frac{dN_{\mathbf{q}}}{dt} = \frac{2\pi}{\hbar}\int_{\varphi=0}^{2\pi}\int_{\theta=0}^{\pi}\int_{0}^{\pi/d}\big\{|\langle \mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 \times f_e(\mathbf{k}+\mathbf{q}) - |\langle \mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2 \cdot f_e(\mathbf{k})\big\} \times \delta\big(E_e(\mathbf{k}+\mathbf{q}) - E_e(\mathbf{k}) - \hbar\omega_{\mathrm{LO}}\big)\frac{2k^2V\sin\theta dkd\theta d\varphi}{(2\pi)^3}. \tag{11}$$

Let's direct the axis OZ in the direction of the vector $\mathbf{q}$. It should be noted that the integral expression does not depend on the angle φ, but $f_e(\mathbf{k}+\mathbf{q}) = f_e(k, q, \cos\theta)$.

Let us denote the argument of the δ-function in (11) by $I(\cos\theta)$ and express the energies $E_e(\mathbf{k}+\mathbf{q})$ and $E_e(\mathbf{k})$ using

the law of dispersion (6):

$$I(\cos\theta) = \frac{1}{2\alpha}\left(\sqrt{1 + \frac{2\alpha\hbar^2(\mathbf{k}^2 + \mathbf{q}^2 + 2qk\cdot\cos\theta)}{m_e}} - 1\right) - \frac{1}{2}\left(\sqrt{1 + \frac{2\alpha\hbar^2\mathbf{k}^2}{m_e}} - 1\right) - \hbar\omega_{\mathrm{LO}}. \quad (12)$$

Let us replace the argument of the δ-function in (11) with $\cos\theta$:

$$\delta(\cos\theta) = \frac{\delta(\cos\theta - \cos\theta_{\mathrm{root}})}{\left|\frac{dI}{d\cos\theta}\big|_{\cos\theta=\cos\theta_{\mathrm{root}}}\right|} = \frac{\delta(\cos\theta - \cos\theta_{\mathrm{root}})\sqrt{1 + \frac{2\alpha\hbar^2(k^2+q^2+2qk\cdot\cos\theta_{\mathrm{root}})}{m_e}}}{\hbar^2 qk}. \quad (13)$$

The root of the argument of δ-function in (11) is

$$\cos\theta_{\mathrm{root}} = \frac{\left(\left(\sqrt{1+\frac{2\alpha\hbar^2k^2}{m_e}} + 2\alpha\hbar\omega_{\mathrm{LO}}\right)^2 - 1\right)\frac{m_e}{2\alpha\hbar^2} - q^2 - k^2}{2qk}. \quad (14)$$

Thus, the expression for the rate of change in the number of phonons can be transformed as follows:

$$\frac{dN_q}{dt} = \frac{4\pi^2}{\hbar}\int\limits_{\theta=0}^{\pi}\int\limits_{0}^{\pi/d}\{|\langle\mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 \times f_e(k, q, \cos\theta) - |\langle\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2\cdot f_e(\mathbf{k})\} \times \frac{\delta(\cos\theta - \cos\theta_{\mathrm{root}})\left(\sqrt{1+\frac{2\alpha\hbar^2k^2}{m_e}} + 2\alpha\hbar\omega_{\mathrm{LO}}\right)}{\hbar^2 q} \times \frac{2kV\,dk\,d\cos\theta}{(2\cdot\pi)^3} = \frac{V}{\pi\hbar}\int\limits_{k\,\min}^{k\,\max}\{|\langle\mathbf{k}, N_q+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2 \times f_e(k, q, \cos\theta_{\mathrm{root}}) - |\langle\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2\cdot f_e(\mathbf{k})\} \times \frac{\left(\sqrt{1+\frac{2\alpha\hbar^2k^2}{m_e}} + 2\alpha\hbar\omega_{\mathrm{LO}}\right)}{\hbar^2 q}kdk. \quad (15)$$

It should be noted that $f_e(k, q, \cos\theta_{\mathrm{root}}) = f_e(\mathbf{k}) \times \exp\left(-\frac{\hbar\omega_{\mathrm{LO}}}{k_{\mathrm{B}}T_e}\right)$.

The integration interval over k should satisfy the law of conservation of energy. This law is implemented in the region (see Figure 1), where the modulus of the phonon wave vector q ensures non-negativity k for any allowable angle θ between them. This area is bounded from below by the intersection of the roots k_{root} of the function $I(\cos\theta)$ at the boundary values $\theta = 0$ and $\theta = \pi$. As a result

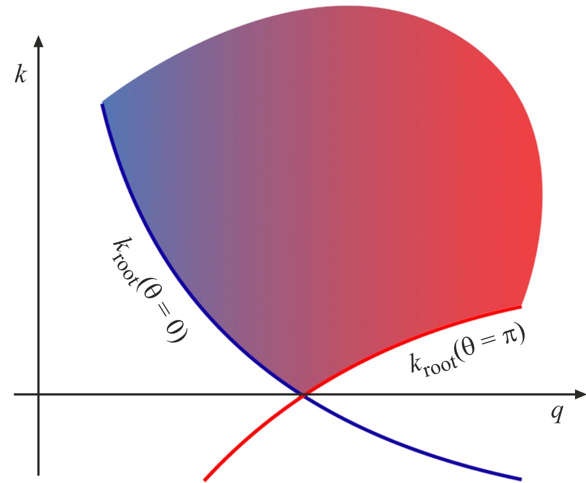

Figure 1. The range of acceptable values of k and q, satisfying the law of conservation of energy at any allowable polar angle, is highlighted with a colored gradient. The blue and red solid curves show the dependencies of k_{root} at the boundary values of $\theta = 0$ and $\theta = \pi$, respectively. (A color version of the figure is provided in the online version of the paper.)

$$\frac{dN_q}{dt} = \frac{V}{\pi\hbar^3 q}\int\limits_{|k_{\mathrm{root}}(\theta=0)|}^{\pi/d}\left\{|\langle\mathbf{k}, N_{\mathbf{q}}+1|H_{\mathrm{LO}}|\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}\rangle|^2\cdot f_e(\mathbf{k}) \times \exp\left(-\frac{\hbar\omega_{\mathrm{LO}}}{k_{\mathrm{B}}T_e}\right) - |\langle\mathbf{k}+\mathbf{q}, N_{\mathbf{q}}-1|H_{\mathrm{LO}}|\mathbf{k}, N_{\mathbf{q}}\rangle|^2\cdot f_e(\mathbf{k})\right\} \times \left(\sqrt{1+\frac{2\alpha\hbar^2k^2}{m_e}} + 2\alpha\hbar\omega_{\mathrm{LO}}\right)kdk. \quad (16)$$

Using expressions for squares of matrix elements (7) and (8) and the nonequilibrium DF of phonons (3), we rewrite the expression for $(dN_{\mathbf{q}})/dt$, getting rid of the recursion:

$$\frac{dN_{\mathbf{q}}}{dt} = \frac{\left((N_{\mathbf{q}0}+1)\cdot\exp\left(-\frac{\hbar\omega_{\mathrm{LO}}}{k_{\mathrm{B}}T_e}\right) - N_{\mathbf{q}0}\right)\beta}{1 - \tau_q^{\mathrm{LO}}\left(\exp\left(-\frac{\hbar\omega_{\mathrm{LO}}}{k_{\mathrm{B}}T_e}\right) - 1\right)\beta}, \quad (17)$$

where

$$\beta = A\int\limits_{|k_{\mathrm{root}}(\theta=0)|}^{\pi/d}\exp\left(-\frac{E_e(k)}{k_{\mathrm{B}}T_e}\right) \times \left(\sqrt{1+\frac{2\alpha\hbar^2\mathbf{k}^2}{m_e}} + 2\alpha\hbar\omega_{\mathrm{LO}}\right)kdk, \quad (18)$$

$$A = \frac{1}{\pi\hbar^3}\frac{2\pi\hbar^2 eE_0}{m_e q^3}\left(1 + (qr_{\mathrm{D}})^{-2}\right)^{-2}\frac{n}{N_c}. \quad (19)$$

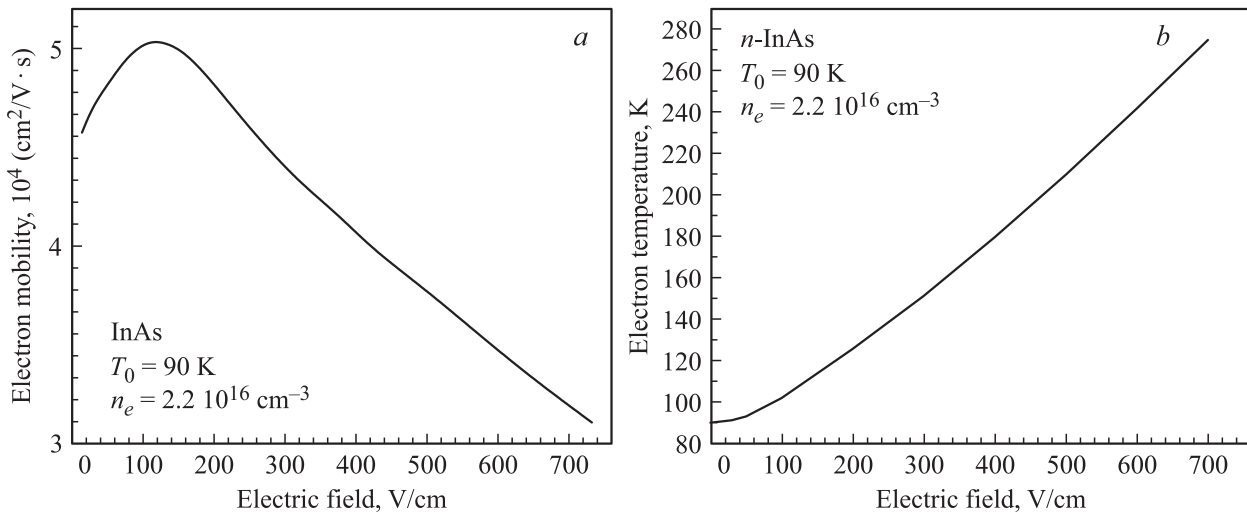

Figure 2. Field dependence of electron mobility (according to experimental data from [17]) (*a*). The results of calculation of the dependence of the electron temperature on the electric field using the power balance equation in the form (20) (*b*).

Using (4), we write the power balance equation (1) as follows:

$$e\mu_e \mathbf{F}^2 = \frac{4\pi}{(2\pi)^3 n} \int_0^{\pi/d} \hbar\omega_{\mathrm{LO}} \frac{dN_q}{dt} q^2 dq. \quad (20)$$

The dependence of mobility $\mu_e(F)$ (see Figure 2, *a*) was found from the analysis of the IV characteristic of *n*-InAs sample obtained in Ref. [17]. The dependence of the electron temperature on the electric field is found by substituting the found dependence $\mu_e(F)$ into the balance equation (20) and performing a numerical solution (see Figure 2, *b*).

3. Calculation of the degree of anisotropy of spontaneous emission

When nonequilibrium charge carriers in InAs are excited by radiation with a photon energy of $\hbar\omega$ exceeding the band gap of E_g, luminescence is observed associated with the radiative recombination of nonequilibrium electrons and holes.

The probability of transition per unit of time between the initial $|i\rangle$ and the final $|f\rangle$ states of an electron is determined by Fermi's golden rule in the first order of perturbation theory [23,24]:

$$P_{i\to f} = \frac{2\pi}{\hbar} |\langle f|H_\omega|i\rangle|^2 \delta(E_f - E_i \mp \hbar\omega), \quad (21)$$

where H_ω is the operator of interaction of a quantum system with electromagnetic radiation. Dirac's δ-function reflects the law of conservation of energy. The term $\hbar\omega$ is taken with a minus if the system increases energy by absorbing a photon with energy $\hbar\omega$, and with a plus if the system loses energy to emit radiation.

A schematic representation of the InAs energy bands and interband optical transitions is shown in Figure 3. The energies of electrons and holes are calculated from the bottom of the corresponding bands.

The spectral radiation density during recombination of electrons from the conduction band and holes from the subbands of heavy and light holes in the valence band, taking into account the charge carrier distribution functions is defined in a semiconductor crystal as

$$L_{c\to hh} = \int_{\mathrm{ZB}} \frac{2\pi}{\hbar} |\langle hh|H_\omega|c\rangle|^2 f_e(\mathbf{k}) f_{hh}(\mathbf{k}) \times \delta(E_{hh}(k) + E_e(k) + E_g - \hbar\omega)\rho_\omega \frac{2d\mathbf{k}}{(2\cdot\pi)^3}, \quad (22)$$

$$L_{c\to lh} = \int_{\mathrm{ZB}} \frac{2\pi}{\hbar} |\langle lh|H_\omega|c\rangle|^2 f_e(\mathbf{k}) f_{lh}(\mathbf{k}) \times \delta(E_{lh}(k) + E_e(k) + E_g - \hbar\omega)\rho_\omega \frac{2d\mathbf{k}}{(2\cdot\pi)^3}, \quad (23)$$

where the indices *hh* and *lh* refer to the subbands of heavy and light holes, respectively; $f_e(\mathbf{k})$, $f_{hh}(\mathbf{k})$, $f_{lh}(\mathbf{k})$ — distribution functions of electrons, heavy and light holes describing nonequilibrium charge carriers under conditions of interband photoexcitation and exposure to a homogeneous electric field; $E_e(k)$, $E_{hh}(k)$, $E_{lh}(k)$ — dispersion dependences of the energy of electrons, heavy and light holes; $\rho_\omega = \frac{\omega^2 n_\omega^3}{2\pi^2 c^3}$ — density of photonic states, n_ω — frequency-dependent refractive index, c — speed of light in vacuum.

The external electric field leads to heating and drift of free charge carriers, and the distribution function of charge carriers over states changes. As a result of electron drift in an electric field, anisotropy of the DF occurs in momentum

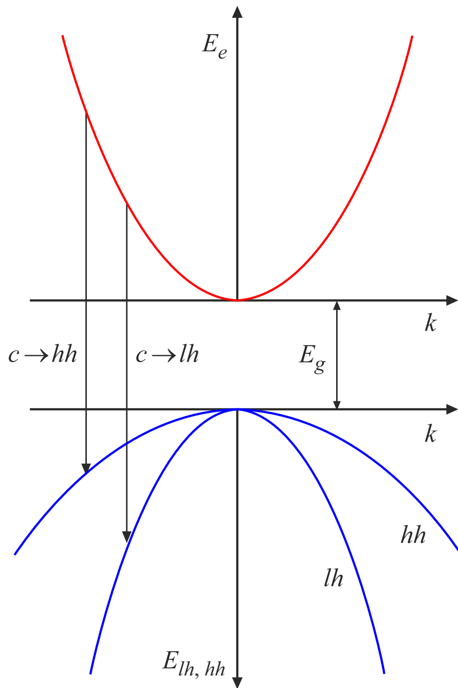

Figure 3. InAs band structure (schematically). Arrows $c \to hh$ and $c \to lh$ show the radiative transitions of electrons from the conduction band to the subbands of heavy and light holes, respectively.

space (a shift of the DF in the direction of electron drift). Then it follows from (22, 23) that, due to the dependence of the optical matrix element of the interband transition on the angle between the electron wave vector and the radiation polarization vector, the intensities of optical transitions for radiation polarized parallel and perpendicular to the drift velocity will differ. Since for the considered concentration of free electrons, the rate of exchange of momentum and energy between charge carriers significantly exceeds the rate of exchange between carriers and scattering centers, the DF can be considered as a Maxwell–Boltzmann function, shifted in velocity space by the drift velocity $\mathbf{v}_{dr}$:

$$f_e(\mathbf{k}) = \frac{n}{N_c}\exp\left(-\frac{E_e(k) - \hbar\mathbf{k}\mathbf{v}_{dr}}{k_B T_e}\right) = \frac{n}{N_c}\exp\left(-\frac{E_e(k) + \hbar\mu_e\mathbf{k}\mathbf{F}}{k_B T_e}\right), \quad (24)$$

where $\mathbf{v}_{dr} = -\mu_e\mathbf{F}$. It should be noted that the change in the concentration of free electrons can be neglected at low levels of interband excitation of nonequilibrium charge carriers compared to the equilibrium case.

We will consider the laws of dispersion of heavy and light holes in InAs in the isotropic and parabolic approximation. The mobility of heavy holes is 2 orders of magnitude lower than the mobility of electrons [25], which is why the DF shift of heavy holes can be neglected. We will also not take into account the anisotropy of the light hole distribution function due to the relatively small contribution of transitions involving light holes to the photoluminescence intensity. Then

$$f_{hh}(\mathbf{k}) = \frac{p}{N_v}\exp\left(-\frac{\hbar^2k^2}{2m_{hh}k_B T_{hh}}\right), \quad (25)$$

$$f_{lh}(\mathbf{k}) = \frac{p}{N_v}\exp\left(-\frac{\hbar^2k^2}{2m_{lh}k_B T_{lh}}\right), \quad (26)$$

where m_{hh}, m_{lh} — effective masses of heavy and light holes, p — concentration of free holes, T_{hh}, T_{lh} — hole temperatures,

$$N_v = \int\limits_{\text{ZB}}\left(\exp\left(-\frac{\hbar^2k^2}{2m_{hh}k_B T_{hh}}\right) + \exp\left(-\frac{\hbar^2k^2}{2m_{lh}k_B T_{lh}}\right)\right)\frac{2dk}{(2\pi)^3}$$

— effective density of states in the valence band.

Let's proceed to integration in spherical coordinates in (29), (30) ($0 \le \theta \le \pi$ — polar angle, $0 \le \varphi \le 2\pi$ — azimuthal):

$$L_{c\to hh} = \int\limits_{\varphi=0}^{2\pi}\int\limits_{\theta=0}^{\pi}\int\limits_{0}^{\pi/d}\frac{2\pi}{\hbar}|\langle hh|H_\omega|c\rangle|^2 \times f_e(\mathbf{k})|_{k=k_{ehh}(\omega)} f_{hh}(\mathbf{k})|_{k=k_{ehh}\omega} \times \frac{\delta(k - k_{ehh}(\omega))}{\left|\frac{d}{dk}(E_{hh}(k) + E_e(k))|_{k=k_{ehh}(\omega)}\right|}\rho_\omega\frac{2k^2\sin\theta dk d\theta d\varphi}{(2\pi)^3}, \quad (27)$$

$$L_{c\to lh} = \int\limits_{\varphi=0}^{2\pi}\int\limits_{\theta=0}^{\pi}\int\limits_{0}^{\pi/d}\frac{2\pi}{\hbar}|\langle lh|H_\omega|c\rangle|^2 \times f_e(\mathbf{k})|_{k=k_{elh}(\omega)} f_{lh}(\mathbf{k})|_{k=k_{elh}\omega} \times \frac{\delta(k - k_{elh}(\omega))}{\left|\frac{d}{dk}(E_{lh}(k) + E_e(k))|_{k=k_{elh}(\omega)}\right|}\rho_\omega\frac{2k^2\sin\theta dk d\theta d\varphi}{(2\pi)^3}, \quad (28)$$

where the values $k_{ehh,elh}(\omega)$ are the roots of the equations $E_{hh,lh}(k) + E_e(k) + E_g - \hbar\omega = 0$.

The optical matrix elements $\langle hh|H_\omega|c\rangle$ and $\langle lh|H_\omega|c\rangle$ depend on the relative position of the polarization vector of the light wave $\mathbf{e}^\omega$ and the vector of the electric field strength $\mathbf{F}$. For certainty, let's choose the axis OZ in the direction of the electric field $\mathbf{F}$ (see Figure 4).

Then, according to Ref. [15], the squares of the matrix elements of interband optical transitions involving heavy holes with perpendicular $|\langle hh|H_\omega|c\rangle_\perp|^2$ and collinear$|\langle hh|H_\omega|c\rangle_\parallel|^2$ vectors $\mathbf{e}^\omega$ and $\mathbf{F}$ are defined as

$$|\langle hh|H_\omega|c\rangle_\parallel|^2 = A\sin^2\theta, \quad (29)$$

$$|\langle hh|H_\omega|c\rangle_\perp|^2 = A\left(\cos^2\theta\cos^2\varphi + \sin^2\varphi\right), \quad (30)$$

where A is a multiplier independent of the values used for integration ($A = \text{const}(k, \theta, \varphi)$). The squares of the matrix

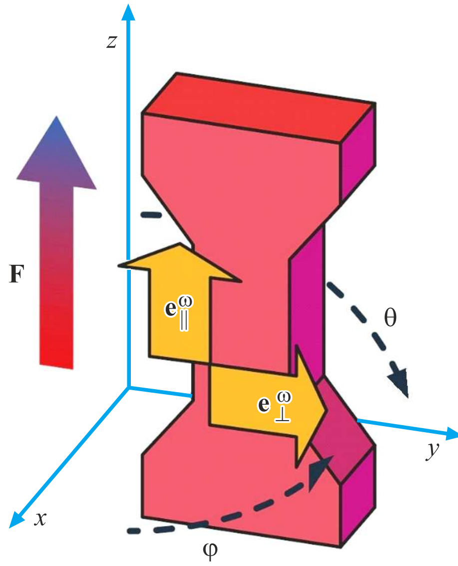

Figure 4. Schematic representation of the sample; $\mathbf{e}_{\parallel}^{\omega}$ and $\mathbf{e}_{\perp}^{\omega}$ — polarization vectors of the light wave directed parallel and perpendicular to the electric field vector **F**.

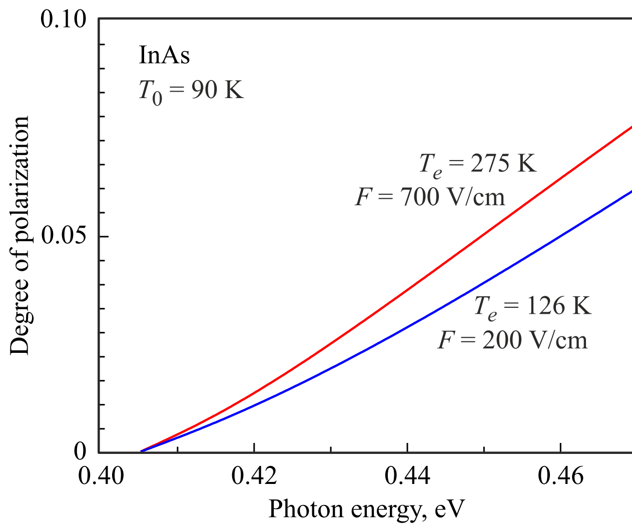

Figure 5. Spectral dependence of the degree of polarization of luminescence radiation in n-InAs at electric field magnitudes of 200 V/cm (blue curve) and 700 V/cm (red curve).

elements of the interband junctions with the participation of light holes have the following form:

$$|\langle lh|H_\omega|c\rangle_{\parallel}|^2 = \frac{A}{3}\left(4\cos^2\theta + \sin^2\theta\right), \quad (31)$$

$$|\langle lh|H_\omega|c\rangle_{\perp}|^2 = \frac{A}{3}\left[\cos^2\varphi\left(4\sin^2\theta + \cos^2\theta\right) + \sin^2\varphi\right], \quad (32)$$

As a result, it is possible to calculate the degree of linear polarization of radiation with the participation of subbands of light and heavy holes of the valence band, depending on the photon energy:

$$P(\omega) = \frac{(L^{\perp}_{c\to hh} + L^{\perp}_{c\to lh}) - (L^{\parallel}_{c\to hh} + L^{\parallel}_{c\to lh})}{(L^{\perp}_{c\to hh} + L^{\perp}_{c\to lh}) + (L^{\parallel}_{c\to hh} + L^{\parallel}_{c\to lh})}. \quad (33)$$

The calculated spectral dependence of the degree of polarization of luminescence radiation under conditions of interband photoexcitation and exposure to a homogeneous electric field is shown in Figure 5. It can be seen that the radiation is dominated by polarization perpendicular to the applied field. With an increase of the photon energy and (or) the electric field, the degree of polarization of luminescence in n-InAs increases.

The calculations used a non-parabolic dispersion law (6) and the electron state distribution function shifted in velocity space (24). As a result, the degree of polarization of photoluminescence in an electric field (33) was calculated numerically. It is possible to qualitatively explain the obtained dependencies shown in Figure 5 using the decomposition of $P(\omega)$ in a series according to the parameter $\frac{m_{ehh}v_{dr}^2}{k_B T_e}$:

$$P(\omega) \propto \frac{m_{ehh}v_{dr}^2}{k_B T_e}\frac{(\hbar\omega - E_g)}{k_B T_e}, \quad (34)$$

where m_{ehh} is the reduced effective mass of electrons and heavy holes. The expression (34) was obtained by studying the anisotropy of photoluminescence polarization in n-GaAs in an electric field for $c \to hh$ transitions in the parabolic approximation [14]. The ratio $\frac{m_{ehh}v_{dr}^2}{k_B T_e}$, proportional to the square of the ratio of the drift and thermal velocities, determines the degree of anisotropy of the distribution function in k-space. It is the anisotropy of the distribution function that results in the polarization dependence of photoluminescence. The deviation of the spectral dependence of the degree of anisotropy from the linear law (34) with an increase in $\hbar\omega$ can be explained by an increase of the effective mass in the conduction band, which depends on the electron energy.

4. Conclusion

The spectral dependence of the degree of linear polarization of luminescence radiation under conditions of interband photoexcitation and application of a homogeneous electric field in a bulk semiconductor n-InAs is calculated. The model took into account the recombination of electrons with both heavy and light holes, and the nonparabolicity of the conduction band was taken into account using the Kane's law of dispersion. When determining the electron temperature, the effect of accumulation of nonequilibrium optical phonons was taken into account.

It is shown that the application of an electric field leads to the appearance of anisotropy of radiation polarization. The luminescence is dominated by radiation polarized in a plane perpendicular to the direction of the applied field.

Funding

This study was supported by grant from the Russian Science Foundation No. 23-12-00036, https://rscf.ru/project/23-12-00036/.

Conflict of interest

The authors declare that they have no conflict of interest.

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Translated by A.Akhtyamov