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Photogalvanic Effect in Semiparabolic Quantum Well

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In this work we studied the charge carriers behaviour in quantum structures where the symmetry with respect to space coordinates and time-reversal symmetry are broken simultaneously. As the model of such structures we considered *finite* semiparabolic quantum well (we considered earlier the case of triangular QW) placed in external magnetic field. We have shown by numerical analysis that the energy spectra of charge carriers in such structures are anisotropic with respect to in-plane (transverse) motion $\epsilon_n(+k_x) \neq \epsilon_n(-k_x)$. This leads to the anisotropy of charge carriers in-plane momentum transfer which, in its turn leads to the anisotropy of photoconductivity $\sigma(+k_x) \neq \sigma(-k_x)$ and as it follows from our calculations, the effect though not very great, could be measurable for the magnetic field of about few T.

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1. Introduction

Several propositions to observe a number of new effects in asymmetric quantum structures were recently discussed in scientific literature [1–3]. Thus, in Ref. [1] it was suggested that if the symmetry with respect to the space coordinates and time-reversal symmetry are broken simultaneously in the nanostructures, then some photovoltaic and magnetoelectric effects have to occur in them. As an example of such structure, in Ref. [1] the asymmetric double-quantum well (QW) in an external magnetic field was considered. In this model, the space asymmetry is introduced by the δ-barriers of different heights, while external magnetic field parallel to the layers provides time-invariance breaking. Later on in [4], the photogalvanic effect in an asymmetric undoped system of three GaAs/AlGaAs quantum wells was studied experimentally and the predictions by [1] concerning this effect were in general confirmed. In Ref. [5], the field-asymmetric transverse magnetoresistance in non-magnetic quantum structures was studied experimentally. In Ref. [6], the anisotropy of electron momentum transfer occurring in the *infinite* triangular quantum well in an external magnetic field due to electron–phonon interaction was considered theoretically. Despite the remarkable insights presented in Refs. [1, 6], the models discussed there should be emended, however, since the model of δ-barriers, as wells as the infinite triangular quantum well are not very realistic and having in mind possible applications of the effect, there is need to examine more realistic ones. The model of finite triangular QW was considered previously in our work [7]. Here we report only the results obtained for the finite semiparabolic QW.

2. Finite semiparabolic quantum well in an external magnetic field

Notice that the photogalvanic effect described above is universal in the sense that it could be observed in QW of different shapes, if only the QW-potential is asymmetric with respect to the space-coordinates inversion. Therefore, consider the semiparabolic QW. It is worthy to mention that the fabrication of semi-parabolic QWs were reported in the paper [8], while some of their non-linear optical properties were examined in [2, 3]. To treat the two-dimensional electron gas in an external magnetic field, we start with conventional approach based on effective mass equation of the form [9]:

$$\left[E_{\mathrm{c}} + \frac{(\mathrm{i}\hbar\nabla + e\boldsymbol{A})^2}{2m^*} + U(z)\right]\psi(\boldsymbol{r}, z) = E\psi(\boldsymbol{r}, z), \quad (1)$$

where E_{c} stands for the bottom of the semiconductor conduction band, $\boldsymbol{A}$ — vector potential, e and m^* are the electron charge and effective mass, respectively, $\boldsymbol{r} = (x, y)$ is in-plane two-dimensional vector. Let us take the next gauge: $\boldsymbol{A} = (Bz, 0, 0)$, then in case of semiparabolic QW, Eq. (1) reduces to the next one

$$\left[\frac{\hat{p}_z^2}{2m^*} + \frac{(eBz + \hbar k_x)^2}{2m^*} + \alpha z^2 - \varepsilon\right]\varphi(k_x, z) = 0. \quad (2)$$

Here $0 \leq z \leq d$, d is the QW-width while the coefficient α can be defined as follows: $\alpha = \frac{1}{2}m^*\omega_0^2 = U_0/d^2$ and $\omega_0^2 = 2U_0/d^2 m^*$, where U_0 is the QW-depth. Introduce also two additional auxiliary parameters, $\omega_{\mathrm{c}} = |e|B/m^*$, $\omega_{\mathrm{c0}}^2 = \omega_{\mathrm{c}}^2 + \omega_0^2$; then instead of Eq. (2) we arrive at the following effective mass equation:

$$\left[\frac{\hat{p}_z^2}{2m^*} + \frac{1}{2}m^*\omega_{\mathrm{c0}}^2\left(z + \frac{\omega_{\mathrm{c}}^2}{\omega_{\mathrm{c0}}^2}z_k\right)^2 - \varepsilon_1\right]\varphi(k_x, z) = 0,$$

where $\varepsilon_1 = \varepsilon - \frac{1}{2}m^*(\omega_0^2\omega_{\mathrm{c}}^2/\omega_{\mathrm{c0}}^2)z_k^2$, $z_k = \hbar k_x/|e|B$. Introducing new variable $\zeta = \sqrt{m^*\omega_{\mathrm{c0}}/\hbar}z$, one gets the following equation to be solved numerically:

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$$\left[\frac{\mathrm{d}^2}{\mathrm{d}\zeta^2} - \left(\zeta + \frac{\omega_{\mathrm{c}}^2}{\omega_{\mathrm{c}0}^2}\zeta_k\right)^2 + \tilde{\varepsilon}\right]\tilde{\varphi}(\zeta,\zeta_k) = 0.$$

Here $\zeta_k = \sqrt{m^*\omega_{\mathrm{co}}/\hbar}z_k$, $\tilde{\varepsilon} = 2\varepsilon_1/\hbar\omega_{\mathrm{c}0}$, and $0 \le \zeta \le \zeta_0 = \sqrt{m^*\omega_{\mathrm{c}0}/\hbar}d$. The results of numerical solution of the last equation for some values of magnetic field B and the boundary conditions $\phi(0) = 0, \phi(d)$ are shown in Table.

TABLE I

Results of numerical solution of the last equation.

$U_{0(\mathrm{e})} = 0.207$ eV, $U_{0(\mathrm{h})} = 0.138$ eV						
	$d = 10$ nm	$d = 15$ nm	$d = 20$ nm	$d = 10$ nm	$d = 15$ nm	$d = 20$ nm
	$B = 0.5$ T			$B = 2.5$ T		
$n = 0$	0.0289207	0.0234808	0.0207753	0.0296978	0.0244304	0.0218671
$n = 1$	0.0509498	0.0382208	0.0318787	0.0521366	0.0396803	0.0335664
$n = 2$	0.0729548	0.0529318	0.0429485	0.0744556	0.0547833	0.0450957
$n = 3$	0.0921393	0.0676281	0.0540012	0.0967151	0.0698132	0.0565407

It is worthy to mention that unlike the previous case of finite triangular QW, here we do not use any small parameter.

Using the standard approach (see, for instance [10]), one gets the general expression for the real part of the photoconductivity which we denote by σ:

$$\sigma = \frac{\pi e^2}{m_{\mathrm{e}}^2\omega}\frac{2}{\Omega}\sum_{i,j}|\langle|j|\boldsymbol{e}\cdot\hat{\boldsymbol{p}}|i\rangle|^2[f(E_i) - f(E_j)]$$
$$\times\delta(E_j - E_i - \hbar\omega),$$

where, since the electromagnetic wave is polarized along z, $\boldsymbol{e} = (0,0,1)$ and $\boldsymbol{e}\cdot\hat{\boldsymbol{p}} = -\mathrm{i}\hbar\partial/\partial z$. The factor of 2 in front of the summation is for spin, $f(E_i), f(E_j)$ are the corresponding Fermi-factors and Ω stands for the volume of the system. Then doing in a similar way as in Ref. [10], after some manipulations which include the summation of Dirac-comb, one gets the following expression for the inter-subband transitions in QW:

$$\sigma^{\pm} = \frac{\pi e^2}{m_{\mathrm{e}}^2 d\omega}\sum_{n,m}|\boldsymbol{e}\cdot\boldsymbol{p}_{\mathrm{cn,vm}}|^2|\langle\mathrm{cn}|\mathrm{vm}\rangle|^2$$
$$\times\left(m_{\mathrm{e}}\tilde{m}^{*(\pm)}_{\mathrm{cn,vm}}/\pi\hbar^2\right)\Theta\left[\hbar\omega - (E_{\mathrm{g}} + \epsilon_{\mathrm{cn}} - \epsilon_{\mathrm{vm}})\right]. \quad (3)$$

Here

$$\boldsymbol{e}\cdot\boldsymbol{p}_{\mathrm{cn,vm}}\langle\mathrm{cn}|\mathrm{vm}\rangle$$
$$\equiv \boldsymbol{e}\cdot\boldsymbol{p}_{\mathrm{cn,vm}}\int\varphi^*_{\mathrm{cn}}(z)\varphi_{\mathrm{vm}}\,\mathrm{d}z \approx \langle\mathrm{cn}\boldsymbol{k}|\boldsymbol{e}\cdot\hat{\boldsymbol{p}}|\mathrm{vm}\boldsymbol{k}\rangle,$$

where "c" and "v" stand for the conduction and valence bands, n and m enumerate the bound states within the QWs and $\boldsymbol{k}$ stands for the in-plane (transverse) wave vector, while the matrix element $\boldsymbol{e}\cdot\boldsymbol{p}_{\mathrm{cn,vm}}$ depends on the nature of Bloch functions and on the polarization $\boldsymbol{e}$; $\Theta(x)$ is the step function and the $\pm$-superscripts of σ and $\tilde{m}^*_{cn,vm}$ correspond to $+k_x$ and $-k_x$, respectively. The main difference between formula (3) and the analogous formula from Ref. [10], is that here, instead of reduced effective mass m^*_{cv} we have $\tilde{m}^{*\pm}_{cn,mv}$, which are defined as follows:

$$\left(\tilde{m}^{*+}_{\mathrm{cn,vm}}\right)^{-1} = \left(\tilde{m}^{*+}_{\mathrm{cn}}\right)^{-1} + \left(\tilde{m}^{*+}_{\mathrm{vm}}\right)^{-1},$$

$$\left(\tilde{m}^{*-}_{\mathrm{cn,vm}}\right)^{-1} = \left(\tilde{m}^{*-}_{\mathrm{cn}}\right)^{-1} + \left(\tilde{m}^{*-}_{\mathrm{vm}}\right)^{-1}.$$

Here $\tilde{m}^{\pm}_{\mathrm{c}n}$ and $\tilde{m}^{\pm}_{\mathrm{v}m}$ are the renormalized effective masses of charge carriers in the QWs of conduction and valence bands, respectively. The dispersion curve for the subband $n = 1$ is shown in Fig. 1, while the results of our calculations for the anisotropy of the photoconductivity are presented in Fig. 2.

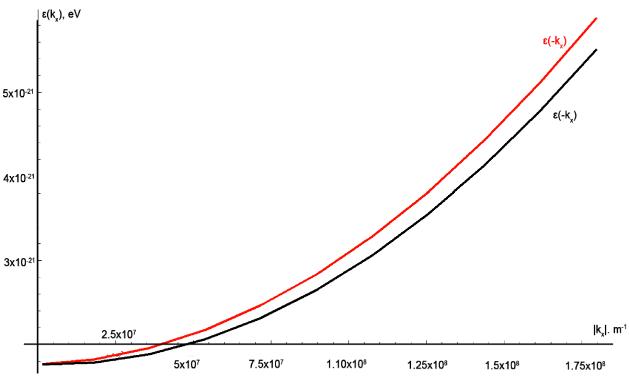

Fig. 1. The dispersion curve.

It is clearly seen from Fig. 2 that the anisotropy of the photoconductivity should occur in such asymmetric nanostructures in an external magnetic field and that the effect should be measurable, since $\Delta\sigma/\sigma(0)$-ratio is about 0.016 at the moderate magnetic field of about ≈ 5 T.

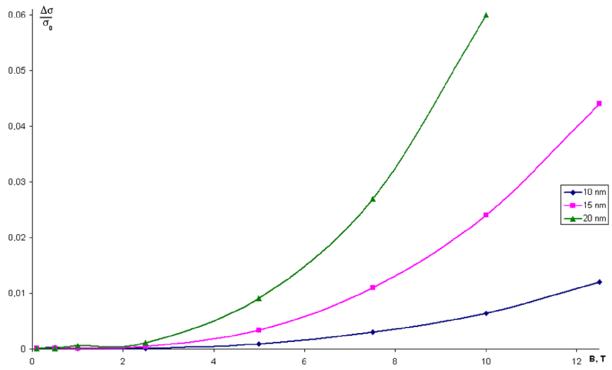

Fig. 2. The anisotropy of the photoconductivity.

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