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Photogalvanic current in a double quantum well

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A double quantum well affected by external alternating electric field with in- and out-of-plane components is studied. This field causes transitions between near-degenerate states located in different wells. The phototransitions are accompanied by the in-plane momentum non-conservation caused by the impurity scattering. We study the in-plane stationary current due to the lack of the in-plane symmetry of these indirect phototransitions. It is shown that the value and direction of the current are determined by the polarization of light. The linear and circular photogalvanic coefficients are found. When the photon energy approaches the distance between subbands these coefficients have their symmetric and antisymmetric resonance behavior, correspondingly.

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1. Introduction. The stationary photocurrent caused by the lack of the system inversion symmetry in homogeneous systems (photogalvanic effect, PGE) was a subject of a wide circle of literature, see reviews [1–6], and the activity in this field continues. In confined systems the photogalvanic current along the surface exists due to non-equivalence of the directions across the surface, even if crystal asymmetry is negligible [7–12]. In this case the current occurs if the electric field of the light has both in- and out-plane components.

A simple model of this effect has been discussed in our recent paper [13]. We have considered the PGE in a classical parabolic well in which electrons experience non-homogeneous (across the well) liquid friction. In this system the difference of electron friction on different well sides causes unequal stalling of the vibrational or rotational motion in the alternating electric field that, in turn, leads to the electron drift along the well.

The phenomenology of PGE in a confined system is determined by the relation for the current density

$$\mathbf{j} = \alpha_s([\mathbf{E} - \mathbf{n}(\mathbf{nE})](\mathbf{nE}^* + \text{c.c}) + i\alpha_a[\mathbf{n}[\mathbf{EE}^*]], \quad (1)$$

where $\mathbf{n}$ is the normal to the quantum well, $\mathbf{E}(t) = \mathrm{Re}(\mathbf{E}e^{-i\omega t})$ is the alternating electric field of light. Real constants α_s and α_a describe linear and circular photogalvanic effects, correspondingly.

The present paper deals with the PGE in a double quantum well, where the light causes the transitions between closely spaced collectivized quantum subbands originating from individual wells. This system looks perspective because the structure of the levels of a double quantum well permits easy tuning of the distance between subbands to the frequency of the external field.

The PGE in a double quantum well is illustrated by Fig. 1. Since the appearance of the photogalvanic current requires non-conservation of the in-plane momentum in the electron excitation process, the phototransitions should include the participation of the "third body". In our case the impurities play the role of this agent.

We consider intersubband transitions of electrons in a system with the quadratic energy spectrum. An electron goes between two states $\epsilon_n(\mathbf{p})$ and $\epsilon_{n'}(\mathbf{p}')$ due to the simultaneous action of electric field and scattering. These states originate from mixing the states of different individual quantum wells. The in-plane current appears due to the change of electron in-plane momentum. To "memorize" electric field in- and out-plane components, the transition probability should contain their product. For non-conservation of the electron momentum the scattering should be taken into account. This transition probability arises in the second order of the perturbation theory. The amplitude of transitions has a resonance on an intermediate state. The subbands of the quantum well are equidistant, that gives rise to the absence of the resonance smearing due to the difference in electron momenta. The result of excitation is the pumping of the momentum to the electron subsystem and the in-plane current.

The paper is organized as follows. First, we will find the transition probability in a classical electric field. Then, the current will be found using a many-band ki-

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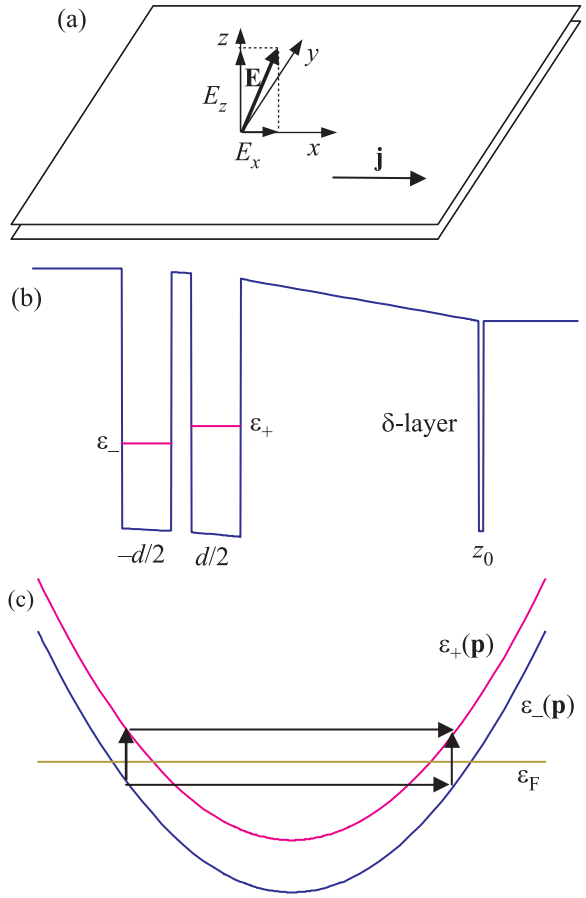

Fig. 1. (a) – The scheme of the proposed experiment. Electric field of light $\mathbf{E}(t)$ has both in-plane and out-of-plane components. (b) – The sketch of the band structure. Quantum wells are centered in planes $z = \pm d/2$. The carriers are provided by the δ-layer of donors in plane z_0. (c) – The transition diagram. The transition amplitude includes vertical transition caused by the light between $\pm$ subbands and impurity scattering which does not conserve the in-plane momentum

netic equation. The last paragraph is devoted to the discussion of the obtained results.

Stationary photocurrent in a double well. We study electrons with a parabolic isotropic energy spectrum in a double quantum well (see Fig. 1). The amplitude of transition between wells is weak, but comparable to the separation of energies of individual wells. The states with in-plane electron momentum $\mathbf{p}$ and subband number $n = \pm$ $|n, \mathbf{p}\rangle = \chi_n(z)\exp(i\mathbf{p}\boldsymbol{\rho})/\sqrt{S}$, ($S$ is the system area, we set $\hbar = 1$) have energies $\epsilon_{n,\mathbf{p}} = p^2/2m + \varepsilon_n$. In this case, the subbands are parallel, $\epsilon_{+,\mathbf{p}} - \epsilon_{-,\mathbf{p}} \equiv \varepsilon_+ - \varepsilon_-$. This circumstance plays an important role in the further consideration, providing the resonance of optical frequency with a distance between subbands for electrons with arbitrary momenta. The overlapping of wave functions $\chi_n(z)$ is supposed to be weak and intersubband distance $\varepsilon_+ - \varepsilon_- = \Delta$ ($\Delta > 0$) is small as compared to the Fermi energy. The scatterers (donors) are distributed in a delta-layer at $z = z_0 > 0$. The well widths and the distance d between them are assumed to be small as compared to z_0.

Assuming that the mean free time is large as compared to the distance between the levels of quantum wells (and also the Fermi energy), one can treat n and $\mathbf{p}$ as good quantum numbers and describe the problem within the kinetic equation for distribution functions $f_{n,\mathbf{p}}$. In such an equation, external classical alternating electric field $\mathbf{E}(t)$ causes the transition between unpertubed states and determines the generation term in the kinetic equation.

The kinetic equation for the stationary distribution function reads

$$\sum_{n',\mathbf{p}'} W^{\rm imp}_{n,\mathbf{p};n',\mathbf{p}'}(f^{(1)}_{n',\mathbf{p}'} - f^{(1)}_{n,\mathbf{p}}) + G_{n,\mathbf{p}} = 0. \qquad (2)$$

Here the first term represents the relaxation due to the scattering processes. Quantity $f^{(1)}_{n,\mathbf{p}}$ is the first correction to equilibrium distribution function $f^{(0)}_{n,\mathbf{p}}$. Function $G_{n,\mathbf{p}}$ is the generation caused by a combined action of the external electric field and the scattering. This term is quadratic in the electric field. Using classical kinetic Eq. (2) means neglecting the off-diagonal elements of the density matrix that is valid if the collision broadening of subbands is less than the distance between them.

Generation $G_{n,\mathbf{p}}$ is given by

$$G_{n,\mathbf{p}} = \sum_{n',\mathbf{p}'} W^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'}(f^{(0)}_{n',\mathbf{p}'} - f^{(0)}_{n,\mathbf{p}}). \qquad (3)$$

We restrict ourselves by the case when the scattering is determined by the charged impurities, so that $W^{\rm imp}_{n,\mathbf{p};n',\mathbf{p}'}$ is attributed to the impurities and $W^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'}$ – to impurities and the electromagnetic field. Then $W^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'}$ is determined by the second order perturbation term which includes the Hamiltonian of the interaction with electromagnetic field $\hat{H}^{\rm ph}$ and the potential energy of the electron interacting with impurities $\hat{V}$. Operator $\hat{H}^{\rm ph}$ is

$$\hat{H}^{\rm ph} = \frac{e}{c}{\rm Re}\left(\mathbf{A}e^{-i\omega t}\right)\hat{\mathbf{v}} \equiv \frac{1}{2}(\hat{U}e^{-i\omega t} + {\rm h.c.}), \qquad (4)$$

where ${\rm Re}(\mathbf{A}e^{-i\omega t})$ is the vector potential of electromagnetic field with frequency ω, $\hat{\mathbf{v}} = (\hat{\mathbf{v}}^{\parallel}, \hat{v}^z)$ is the velocity operator. The complex amplitude of the electric field is $\mathbf{E} = i\omega\mathbf{A}/c$. Thus, the operator $\hat{U} = e(\mathbf{E}\hat{\mathbf{v}})/i\omega$. Note that we suppose the electric field to be homogeneous. The diagonal elements of in-plane components

of the velocity operator $\mathbf{v}^{\parallel}_{n,\mathbf{p};n',\mathbf{p}'} = \mathbf{v}_{\mathbf{p}}\delta_{nn'}\delta_{\mathbf{p},\mathbf{p}'}$, $\mathbf{v}_{\mathbf{p}} = \partial_{\mathbf{p}}\epsilon_{n,\mathbf{p}} = \mathbf{p}/m$. The normal component has matrix elements $v^z_{n,\mathbf{p};n',\mathbf{p}'} = v^z_{n,n'}\delta_{\mathbf{p},\mathbf{p}'}$. The impurity potential reads

$$V(\mathbf{r}) = \sum_i u(\mathbf{r} - \mathbf{r}_i), \tag{5}$$

where the sum runs over all the impurities situated in points $\mathbf{r}_i$ with individual potentials $u(\mathbf{r} - \mathbf{r}_i)$.

The excitation probability including the impurity scattering is determined by the second-order transition amplitude. The needed term arises from the interference of amplitudes caused by the E_z and in-plane components of the electric field. The draft of the transitions is depicted in Fig. 1.

In the second order of the interaction, the transition probability is

$$W^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'} = \frac{\pi}{2}\Biggl\{\delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} + \omega) \times \Biggl\langle\Biggl|\sum_{n_1}\Biggl[\frac{V_{n,\mathbf{p};n_1,\mathbf{p}'}U^+_{n_1,\mathbf{p}';n',\mathbf{p}'}}{\eta + i(\varepsilon_{n_1,n'} + \omega)} + \frac{U^+_{n,\mathbf{p};n_1,\mathbf{p}}V_{n_1,\mathbf{p};n',\mathbf{p}'}}{\eta + i(\varepsilon_{n_1,n} - \omega)}\Biggr]\Biggr|^2\Biggr\rangle + \delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} - \omega) \times \Biggl\langle\Biggl|\sum_{n_1}\Biggl[\frac{V_{n,\mathbf{p};n_1,\mathbf{p}'}U_{n_1,\mathbf{p}';n',\mathbf{p}'}}{\eta + i(\varepsilon_{n_1,n'} - \omega)} + \frac{U_{n,\mathbf{p};n_1,\mathbf{p}}V_{n_1,\mathbf{p};n',\mathbf{p}'}}{\eta + i(\varepsilon_{n_1,n} + \omega)}\Biggr]\Biggr|^2\Biggr\rangle\Biggr\};$$

$$(\eta = +0). \tag{6}$$

Here $\varepsilon_{n_1,n} \equiv \varepsilon_{n_1} - \varepsilon_n$; angular brackets denote the average over impurities configuration.

The denominators in Eq. (6) have their resonance with the field frequency independently from the electron momentum. At the same time, the resonance in the final state is absent due to non-conservation of the in-plane momentum.

Eq. (6) can be rewritten in the form ($\mathbf{E} = (\mathbf{E}_{\parallel}, E_z)$):

$$W^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'} = \frac{\pi e^2}{2\omega^2} \times \Biggl(\Biggl\langle\Biggl|\sum_{n_1}\Biggl\{V_{n,\mathbf{p};n_1,\mathbf{p}'}\Biggl[\frac{\mathbf{v}_{\mathbf{p}'}\mathbf{E}^*_{\parallel}\delta_{n_1,n'}}{i\omega} + \frac{v^z_{n_1,n'}E^*_z}{\eta + i(\varepsilon_{n_1,n'} + \omega)}\Biggr] + \Biggl[\frac{\mathbf{v}_{\mathbf{p}}\mathbf{E}^*_{\parallel}\delta_{n,n_1}}{-i\omega} + \frac{v^z_{n,n_1}E^*_z}{\eta + i(\varepsilon_{n_1,n} - \omega)}\Biggr]V_{n_1,\mathbf{p};n',\mathbf{p}'}\Biggr\}\Biggr|^2\Biggr\rangle \times \delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} + \omega) + \Biggl\langle\Biggl|\sum_{n_1}\Biggl\{V_{n,\mathbf{p};n_1,\mathbf{p}'}\Biggl[\frac{\mathbf{v}_{\mathbf{p}'}\mathbf{E}_{\parallel}\delta_{n_1,n'}}{-i\omega} + \frac{v^z_{n_1,n'}E_z}{\eta + i(\varepsilon_{n_1,n'} - \omega)}\Biggr] + \Biggl[\frac{\mathbf{v}_{\mathbf{p}}\mathbf{E}_{\parallel}\delta_{n,n_1}}{i\omega} + \frac{v^z_{n,n_1}E_z}{\eta + i(\varepsilon_{n_1,n} + \omega)}\Biggr]V_{n_1,\mathbf{p};n',\mathbf{p}'}\Biggr\}\Biggr|^2\Biggr\rangle \times \delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} - \omega)\Biggr). \tag{7}$$

It is evident that the contribution to photogalvanic effect is given not by the total transition probability $W^{\rm ph}$, but only its odd in $\mathbf{p}, \mathbf{p}'$ part. For this part, we have the following expression:

$$\tilde{W}^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'} = \frac{\pi e^2}{m\omega^3}\Biggl[\Biggl\langle{\rm Re}\Biggl(\sum_{n_1}\Biggl\{V_{n,\mathbf{p};n',\mathbf{p}'}(\mathbf{p}' - \mathbf{p})\mathbf{E}^*_{\parallel} \times \Biggl[\frac{V^*_{n,\mathbf{p};n_1,\mathbf{p}'}v^z_{n',n_1}E_z}{i\eta + (\varepsilon_{n_1,n'} + \omega)} + \frac{v^z_{n_1,n}E_zV^*_{n_1,\mathbf{p};n',\mathbf{p}'}}{i\eta + (\varepsilon_{n_1,n} - \omega)}\Biggr]\Biggr\}\Biggr)\Biggr\rangle \times \delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} + \omega) + \Biggl\langle{\rm Re}\Biggl(\sum_{n_1}\Biggl\{V_{n,\mathbf{p};n',\mathbf{p}'}(\mathbf{p} - \mathbf{p}')\mathbf{E}_{\parallel} \times \Biggl[\frac{V^*_{n,\mathbf{p};n_1,\mathbf{p}'}v^z_{n',n_1}E^*_z}{i\eta + (\varepsilon_{n_1,n'} - \omega)} + \frac{v^z_{n_1,n}E^*_zV^*_{n_1,\mathbf{p};n',\mathbf{p}'}}{i\eta + (\varepsilon_{n_1,n} + \omega)}\Biggr]\Biggr\}\Biggr)\Biggr\rangle \times \delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n',\mathbf{p}'} - \omega)\Biggr]. \tag{8}$$

The current density is expressed via the first angular harmonics of $f^{(1)}_{n,\mathbf{p}}$:

$$\mathbf{j} = \frac{2e}{S}\sum_{n,\mathbf{p}}\mathbf{v}_{\mathbf{p}}f^{(1)}_{n,\mathbf{p}}. \tag{9}$$

Owing to this fact, kinetic equation Eq. (2) can be transformed to algebraic form

$$\frac{1}{\tau_n(p)}f^{(1)}_{n,\mathbf{p}} - \frac{1}{\tau_{n,-n}(p)}f^{(1)}_{-n,\mathbf{p}} = G_{n,\mathbf{p}}, \tag{10}$$

where one should keep the first angular harmonic of $G_{n,\mathbf{p}}$ only. In Eq. (10), $\tau_n(p)$ is the intra-subband transport relaxation time and $\tau_{n,-n}(p)$ is the time of transition from state $(n,\mathbf{p})$ to all states of the subband $(-n)$. These values are determined by

$$\frac{1}{\tau_n(p)} = 2\pi\sum_{\mathbf{p}'}\Bigl[\Bigl\langle|V_{n,\mathbf{p};n,\mathbf{p}'}|^2\Bigr\rangle\delta(\epsilon_{n,\mathbf{p}} - \epsilon_{n,\mathbf{p}'})\left(1 - \frac{\mathbf{p}\mathbf{p}'}{p^2}\right) + \Bigl\langle|V_{n,\mathbf{p};-n,\mathbf{p}'}|^2\Bigr\rangle\delta(\epsilon_{n,\mathbf{p}} - \epsilon_{-n,\mathbf{p}'})\Bigr];$$

$$\frac{1}{\tau_{n,-n}(p)} = 2\pi\sum_{\mathbf{p}'}\Bigl\langle|V_{n,\mathbf{p};-n,\mathbf{p}'}|^2\Bigr\rangle\delta(\epsilon_{n,\mathbf{p}} - \epsilon_{-n,\mathbf{p}'})\frac{\mathbf{p}\mathbf{p}'}{p^2}. \tag{11}$$

Solving Eq. (10) we find (argument $\mathbf{p}$ is omitted):

$$f_n^{(1)} = \Big(G_n\tau_n + G_{-n}\frac{\tau_+\tau_-}{\tau_{n,-n}}\Big)\Big(1 - \frac{\tau_+\tau_-}{\tau_{+,-}\tau_{-,+}}\Big)^{-1}. \quad (12)$$

The expressions for $\tau_n(p), \tau_{n,-n}(p)$ and $\tilde{W}^{\rm ph}_{n,\mathbf{p};n',\mathbf{p}'}$ contain the correlators of the form $\Big\langle V_{n,\mathbf{p};n',\mathbf{p}'}V_{m,\mathbf{p};m',\mathbf{p}'}\Big\rangle$. In the case of impurities situated in layer $z = z_0$ ($\mathbf{r}_i = (\boldsymbol{\rho}_i, z_0)$), function $V(\mathbf{r})$ reads

$$V(\mathbf{r}) = \sum_{\mathbf{q},i} u_{\mathbf{q}} e^{-q|z-z_0|}\exp\left[-i\mathbf{q}(\boldsymbol{\rho}-\boldsymbol{\rho}_i)\right], \quad (13)$$

where $u_{\mathbf{q}}$ is the 2D Fourier component of the impurity center potential. For example, for unscreened Coulomb center $u_{\mathbf{q}} = 2\pi e^2/\kappa q S$ (κ is the background dielectric constant), the correlators are given by

$$\Big\langle V_{n,\mathbf{p};n',\mathbf{p}'}V^*_{m,\mathbf{p};m',\mathbf{p}'}\Big\rangle = n_s S\int dz dz' |u_{\mathbf{p}-\mathbf{p}'}|^2 \times$$
$$\times\, e^{-q(2z_0-z-z')}\chi_n(z)\chi_{n'}(z)\chi_m(z')\chi_{m'}(z'). \quad (14)$$

Here n_s is the areal density of scatterers. We suppose that the electron wavelength is larger than d. In this approximation one can find from Eq. (14):

$$\Big\langle V_{n,\mathbf{p};n',\mathbf{p}'}V^*_{m,\mathbf{p};m',\mathbf{p}'}\Big\rangle = n_s S|u_{\mathbf{p}-\mathbf{p}'}|^2 e^{-2qz_0} \times$$
$$\times\Big[\delta_{n,n'}\delta_{m,m'} + q(z_{n,n'}\delta_{m,m'} + z_{m,m'}\delta_{n,n'})\Big]. \quad (15)$$

Matrix elements $z_{nn'}$ should be estimated for specific wave functions. For simplicity, we will use the wave functions of two delta-functional wells in the tight-binding approximation. The seed states with energies $\varepsilon_0 \pm \Delta_0/2$ can be written as

$$\chi_{1,2} = \sqrt{\kappa}e^{-\kappa|z\mp d/2|}. \quad (16)$$

In basis (16) $\chi_+ = (1,\beta)/\sqrt{1+\beta^2}$, $\chi_- = (\beta,-1)/\sqrt{1+\beta^2}$, where β is the mixing amplitude. The corresponding states energies are $\varepsilon_\pm = \varepsilon_0 \pm \Delta/2$, $\Delta = \sqrt{\Delta_0^2 + 4t_0^2}$, where $t_0 \sim \varepsilon_0 e^{-\kappa d}$ is a hopping amplitude between wells. For quantity β we have $\beta = 2t_0/(\Delta + \Delta_0)$. The matrix elements of z are $z_{++} = -z_{--} = d(1-\beta^2)/[2(1+\beta^2)]$, $z_{+-} = \beta d/(1+\beta^2)$.

Inserting Eq. (15) into Eq. (11) we get the expressions for $\tau_+ \approx \tau_- = \tau$ and a small difference $1/\tau_- - 1/\tau_+$:

$$\frac{1}{\tau} = mn_s\int\frac{d\mathbf{q}}{2\pi}|\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}\delta(q^2+2\mathbf{pq})\frac{q^2}{p^2},$$
$$\frac{1}{\tau_-} - \frac{1}{\tau_+} = m(z_{++}-z_{--})n_s \times$$
$$\times\int\frac{d\mathbf{q}}{\pi}|\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}\delta(q^2+2\mathbf{pq})\frac{q^3}{p^2}, \quad (17)$$

where $\tilde{u}_{\mathbf{q}} = Su_{\mathbf{q}}$. From Eq. (15) it is seen that $\tau_{n,-n} \gg \gg \tau_n$ and, so, Eq. (12) can be simplified

$$f_n^{(1)} = G_n\tau_n. \quad (18)$$

Further we will consider the resonance situation when frequency ω is close to Δ. Smallness Δ, as compared to the Fermi energy $\epsilon_{\rm F} = p_{\rm F}^2/2m$ ($p_{\rm F}$ being the Fermi momentum) leads to approximate expressions for $G_+ \approx \approx -G_-$,

$$G_+ = -\frac{n_s e^2\Delta z_{+-}^2}{\pi\omega^2}\frac{\partial f^{(0)}_{\mathbf{p}}}{\partial\mu}\int d\mathbf{q}\,\mathrm{Im}\Big(\frac{\mathbf{q}\cdot\mathbf{E}_\parallel E_z^*}{\Delta-\omega+i\eta}\Big)\times$$
$$\times\, q|\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}\delta(q^2+2\mathbf{pq}). \quad (19)$$

For the current we have from Eqs. (9), (12), (19)

$$\mathbf{j} = \frac{n_s e^3\Delta z_{+-}^2}{4\pi^3 m\omega^2}\times$$
$$\times\int d\mathbf{p}\frac{\partial f^{(0)}_{\mathbf{p}}}{\partial\mu}(\tau_+-\tau_-)\mathrm{Im}\Big(\frac{\mathbf{q}\cdot\mathbf{E}_\parallel E_z^*}{\Delta-\omega+i\eta}\Big)\times$$
$$\times\int d\mathbf{q}|\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}q^3 e^{-2qz_0}\delta(q^2+2\mathbf{pq}). \quad (20)$$

Eq. (20) has a resonant character with the resonance at $\omega = \Delta$. This resonance results from the intermediate state for transition due to the parallelism (equidistance) of subbands. The resonance is smeared due to scattering, e.g., by impurities. To include this smearing, infinitesimal η was replaced by finite relaxation rate $1/\tau$ which can be estimated from mobility. This leads to the current finiteness at the resonance point.

At temperature $T = 0$ the latter expression is simplified, and we obtain the final results for the required values in the model of two δ-like wells:

$$\alpha_s = \alpha_0\frac{1}{(\Delta-\omega)^2\tau^2+1};\quad \alpha_a = \alpha_0\frac{\Delta-\omega}{(\Delta-\omega)^2\tau^2+1};$$
$$\alpha_0 = \frac{2e^3 n_e d\tau\Delta}{m\omega^2}\frac{(\beta^2-1)\beta^2}{(1+\beta^2)^3}F, \quad (21)$$

where $n_e = m\epsilon_{\rm F}/\pi$ is the electron concentration. Besides we introduced a dimensionless quantity $F = d^2\Phi_3^2\Phi_2^{-2}$,

$$\Phi_s = \int_0^{2p_{\rm F}} dq q^s |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}\frac{1}{\sqrt{1-q^2/4p_{\rm F}^2}}. \quad (22)$$

In the specific case of $p_{\rm F}z_0 \gg 1$ Eq. (22) is reduced to

$$\Phi_s = \int_0^\infty dq q^s |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}. \quad (23)$$

If the scattering is determined by the charged nonscreened impurities $F = d^2/4z_0^2$.

Discussion. Let us compare the linear and circular effects. Linear coefficient α_s has a constant sign, while circular coefficient α_a changes the sign at $\Delta = \omega$. Namely, α_s has the δ-like frequency dependence and α_a has the antisymmetric resonance. Note that in the vicinity of the resonance $\alpha_s(\omega)$ and $\alpha_a(\omega)$ have same orders and at $|\omega - \Delta|/\tau \gg 1$ $\alpha_a(\omega) \gg \alpha_s(\omega)$. Coefficients α_s and α_a get the extrema at $\Delta = \omega$ and $\Delta = \omega \pm 1/\tau$, correspondingly. The ratio of these absolute values in these points is $\max(\alpha_s)/\max(|\alpha_a|) = 2$. Similar behavior was marked for interband transitions [9, 10] and in a classical parabolic well [13].

In the experiment it can be convenient to scan Δ_0 by means of the gate voltage V at a fixed frequency. In a gated system, parameter Δ_0 depends on the gate voltage as $(V - V_0)d/D$, where D is the distance to the gate, V_0 is the value of V for which $\Delta_0 = 0$. When Δ_0 runs all the range, resonance $\Delta = \omega$ is passed twice.

The typical dependences of the PGE coefficients on the gate voltage via parameter Δ_0 are depicted in the Fig. 2. We have chosen the parameters for

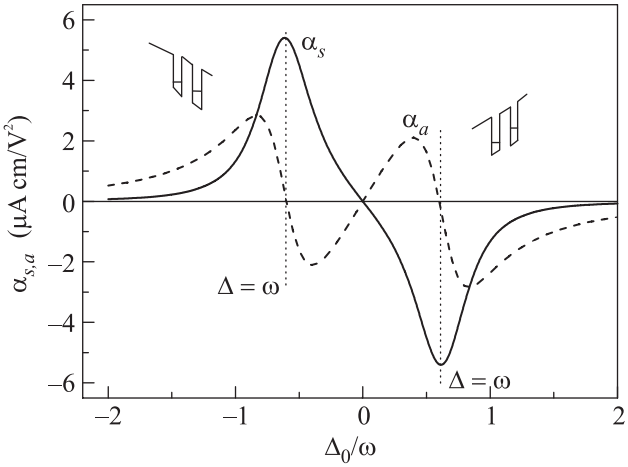

Fig. 2. Dependence of the PGE coefficients α_s (solid) and α_a (dashed) on the bias voltage (expressed via the parameter Δ_0). The dotted lines mark the exact resonance $\omega = \Delta$. The current changes sign with the sign of Δ_0 (see inserts)

the GaAs/AlGaAs double quantum well $d = 10^{-6}$ cm, $z_0 = 3 \cdot 10^{-6}$ cm, $\epsilon_F = 20$ meV ($n_e = 6.2 \cdot 10^{11}$ cm^{-2}), $\Delta = 0.1$ meV, $\tau = 4 \cdot 10^{-11}$ s. Coefficient α_s has a resonance peak and α_a goes through zero when the frequency coincides with the distance between subbands. Besides, both coefficients change the sign at $\Delta_0 = 0$. The found current value looks quite measurable.

It should be emphasized that the initial and final states in the optical transition can belong to different or the same subbands. The resonant behavior results from the resonance on the intermediate state rather than the energy conservation in the final states, because the conservation law for the phototransition with the participation of impurity scattering does not give a fixed frequency for the transition. The sharpness of the resonance is conditioned by the equidistance of the energy bands in a 2D well.

In conclusion, we found the stationary current along a double-well system affected by the linear-polarized far-infrared wave. The stationary current originates from the periodic vibration of electrons between two non-equivalent quantum wells caused by the normal component of the alternating electric field with synchronic in-plane acceleration/deceleration by the in-plane electric field component. The photogalvanic effect needs vertical asymmetry of the double quantum well. The effect has a peak resonant structure connected with the parallel subbands of the double quantum well. The resonant frequency can be easily tuned by the application of the gate voltage. The optimal range of frequencies is 10^{11}–10^{13} s^{-1}. The predicted value of the current is experimentally measurable.

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