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# Pure spin photocurrents in low-dimensional structures

*S. A. Tarasenko*[1], *E. L. Ivchenko*

*A. F. Ioffe Physico-Technical Institute, Russian Academy of Sciences, 194021 St.Petersburg, Russia*

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As is well known the absorption of circularly polarized light in semiconductors results in optical orientation of electron spins and helicity-dependent electric photocurrent, and the absorption of linearly polarized light is accompanied by optical alignment of electron momenta. Here we show that the absorption of unpolarized light leads to generation of a pure spin current, although both the average electron spin and electric current vanish. We demonstrate this for direct interband and intersubband as well as indirect intraband (Drude-like) optical transitions in semiconductor quantum wells.

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**1. Introduction.** Spin and charge are among the basic properties of elementary particles such as an electron, positron and proton. The perturbation of a system of electrons by an electric field or light may lead to a flow of the particles. The typical example is an electric current that represents the directed flow of charge carriers. Usually the electric currents do not entail a considerable spin transfer because of the random orientation of electron spins. However, the charge current can be accompanied by a spin current if electron spins are co-oriented as it happens, e.g., under injection of spin-polarized carriers from magnetic materials [1, 2] or in the optical-orientation-induced circular photogalvanic effect [3, 4]. Furthermore, there exists a possibility to create a pure spin current which is not accompanied by a net charge transfer. This state represents a non-equilibrium distribution when electrons with the spin "up" propagate mainly in one direction and those with the spin "down" propagate in the opposite direction. In terms of the kinetic theory, it can be illustrated by a spin density matrix with two nonzero components, $\rho_{s,s}(\mathbf{k}) = \rho_{\bar{s},\bar{s}}(-\mathbf{k})$, where $s$ and $\mathbf{k}$ are the electron spin index and the wave vector, and $\bar{s}$ means the spin opposite to $s$. Spin currents in semiconductors can be driven by an electric field acting on unpolarized free carriers which undergo a spin-dependent scattering and/or propagate in a medium with spin-orbit coupling. This is the so-called spin Hall effect where a pure spin transfer appears in the direction perpendicular to the electric field, see [5 – 8] and references therein. The spin currents can be induced as well by optical means as a result of interference of one- and two-photon coherent excitation with a two-color electromagnetic field [9] or under interband optical transitions in noncentrosymmetrical semiconductors [10].

Here we show that pure spin currents, accompanied neither by charge transfer nor by spin orientation, can be achieved under absorption of linearly polarized or unpolarized light in semiconductor low-dimensional systems. The effect is considered here for direct interband and intersubband as well as for indirect free-carrier optical transitions in semiconductor quantum wells (QWs).

Phenomenologically, the flux of electron spins is characterized by a pseudotensor $\hat{\mathbf{F}}$ with the components $F_\beta^\alpha$ describing the flow in the $\beta$ direction of spins oriented along $\alpha$, with $\alpha$ and $\beta$ being the Cartesian coordinates. In terms of the kinetic theory such a component of the spin current is contributed by a non-equilibrium correction $\propto \sigma_\alpha k_\beta$ to the electron spin density matrix, where $\sigma_\alpha$ is the Pauli matrix. In general, the concept of spin current is uncertain in systems with spin-orbit interaction, since the spin and spin-dependent velocity cannot be determined simultaneously [11]. Mathematically it is caused by the fact that the Pauli matrices and the velocity operator do not commute. However this problem is overcome in the most of real semiconductor QWs where spin-orbit interaction can be considered as a small perturbation. To the first order in the constant of spin-orbit coupling and within the relaxation time approximation, the pure spin current photoinduced in the conduction band is given by

$$F_\alpha^\beta = \sum_{\mathbf{k}} \tau_e \mathrm{Tr} \left[ \frac{\sigma_\alpha}{2} v_\beta(\mathbf{k}) \, \dot{\rho}(\mathbf{k}) \right] \qquad (1)$$

with the spin-dependent corrections being taken into account either in the velocity operator $\mathbf{v}(\mathbf{k})$ or in the photogeneration rate of the spin density matrix $\dot{\rho}(\mathbf{k})$. Here $\tau_e$ is the relaxation time of the spin current which can differ from the conventional momentum relaxation time that

[1] e-mail: tarasenko@coherent.ioffe.ru

governs the electron mobility. Electron-electron collisions, which do not affect the mobility, can contribute to $\tau_e$ as it happens, e.g., in the case of spin relaxation [12].

**2. Interband optical transitions.** Appearance of a pure spin current under direct optical transitions is linked with two fundamental properties of semiconductor QWs, namely, the linear in the wave vector $\mathbf{k}$ *spin splitting* of energy spectrum and the spin-sensitive *selection rules* for optical transitions [13]. The effect is most easily conceivable for direct transitions between the heavy-hole valence subband $hh1$ and conduction subband $e1$ in QWs of the $\mathrm{C}_s$ point symmetry, e.g. in (113)- or (110)-grown QWs. In such structures the spin component along the QW normal $z$ is coupled with the in-plane electron wave vector. This leads to $\mathbf{k}$-linear spin-orbit splitting of the energy spectrum as sketched in Fig.1, where heavy hole subband $hh1$ is split into two spin

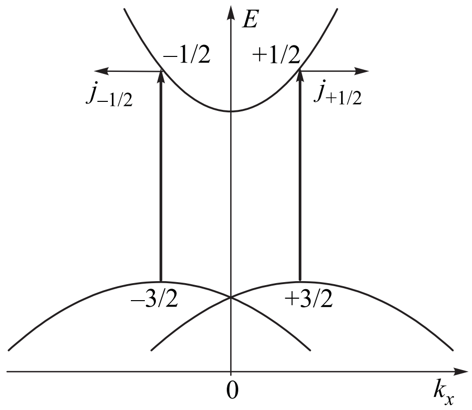


Fig.1. Microscopic origin of pure spin current induced by interband photoexcitation. The vertical lines show the possible optical transitions

branches $\pm 3/2$. As a result they are shifted relative to each other in the $\mathbf{k}$ space. In the reduced-symmetry structures, the spin splitting of the conduction band is usually smaller than that of the valence band and not shown in Fig.1 for simplicity. Due to the selection rules the direct optical transitions from the valence subband $hh1$ to the conduction subband $e1$ can occur only if the electron angular momentum changes by $\pm 1$. It follows then that the allowed transitions are $|+3/2\rangle \rightarrow |+1/2\rangle$ and $|-3/2\rangle \rightarrow |-1/2\rangle$, as illustrated in Fig.1 by vertical lines. Under excitation with linearly polarized or unpolarized light the rates of both transitions are equal. In the presence of spin splitting, the optical transitions induced by photons of the fixed energy $\hbar\omega$ occur in the opposite points of the $\mathbf{k}$ space for the spin branches $\pm 1/2$. The asymmetry of photoexcitation results in a flow of electrons within each spin branch. The corresponding fluxes $\mathbf{j}_{+1/2}$ and $\mathbf{j}_{-1/2}$ are of equal strengths but of opposite directions. Thus, this non-equilibrium electron distribution is characterized by the nonzero spin current $\mathbf{j}_{\rm spin} = (1/2)(\mathbf{j}_{+1/2} - \mathbf{j}_{-1/2})$ but a vanishing charge current, $e(\mathbf{j}_{+1/2} + \mathbf{j}_{-1/2}) = 0$.

The direction $\beta$ of the photoinduced spin current and the orientation $\alpha$ of transmitted spins are determined by the explicit form of spin-orbit interaction. The latter is governed by the QW symmetry and can be varied. For QWs based on zinc blende lattice semiconductors and grown along the crystallographic direction $[110] \parallel z$, the light absorption leads to a flow along $x \parallel [1\bar{1}0]$ of spins oriented along $z$. This component of the electron spin flow can be estimated as

$$F_x^z = \gamma_{zx}^{(hh1)} \frac{\tau_e}{2\hbar} \frac{m_h}{m_e + m_h} \frac{\eta_{cv}}{\hbar\omega} I \,, \qquad (2)$$

where $\gamma_{zx}^{(hh1)}$ is a constant describing the $\mathbf{k}$-linear spin-orbit splitting of the $hh1$ subband, $m_e$ and $m_h$ are the electron and hole effective masses in the QW plane, respectively, $\eta_{cv}$ is the light absorbance, and $I$ is the light intensity.

In (001)-grown QWs the absorption of linearly- or unpolarized light results in a flow of electron spins oriented in the QW plane. In contrast to the low-symmetry QWs considered above, in (001)-QWs the $\mathbf{k}$-linear spin splitting of the $hh1$ valence subband is suppressed and here, for the sake of simplicity, we assume the parabolic spin-independent dispersion in the $hh1$ valence subband and take into account the spin-dependent contribution

$$H_{\rm so}^{(e1)} = \sum_{\alpha\beta} \gamma_{\alpha\beta}^{(e1)} \sigma_\alpha k_\beta \qquad (3)$$

to the electron effective Hamiltonian. Then, to the first order in the spin-orbit coupling, the components of the pure spin current generated in the subband $e1$ are derived to be

$$F_\beta^\alpha = \gamma_{\alpha\beta}^{(e1)} \frac{\tau_e}{2\hbar} \frac{m_e}{m_e + m_h} \frac{\eta_{cv}}{\hbar\omega} I \,. \qquad (4)$$

*Hole spin currents.* For heavy holes $hh1$ one can introduce the pseudospin description with pseudospins $\tilde{s} = \pm 1/2$ representing the hole states with the angular momentum $\pm 3/2$. In addition to the electron spin current (4), a hole pseudospin current is induced by interband photoexcitation with linearly polarized light. In the geometry of normal incidence the components $F_\alpha^x(hh1)$ and $F_\alpha^y(hh1)$ are given, respectively, by the real and imaginary parts of

$$F_\alpha^x + \mathrm{i} F_\alpha^y = (\gamma_{x\alpha}^{(e1)} + \mathrm{i}\gamma_{y\alpha}^{(e1)}) \frac{\tau_h}{2\hbar} \frac{m_e}{m_e + m_h} \frac{\eta_{cv}}{\hbar\omega} I (e_x + \mathrm{i} e_y)^2,$$

where $\tau_h$ is the relaxation time of the hole spin current and $\mathbf{e}$ is the light polarization unit vector.

**3. Intersubband optical transitions.** The intersubband light absorption in $n$-doped QW structures is a resonant process and occurs if the photon energy equals the energy spacing between the subbands. In a simple one-band model, direct optical transitions from the subband $e1$ to the subband $e2$ can be induced only by irradiation with nonzero normal component $e_z$ of the polarization vector. These transitions occur with the spin conservation, $(e1,+1/2) \to (e2,+1/2)$ and $(e1,-1/2) \to (e2,-1/2)$, as illustrated in Fig.2 by vertical solid lines.

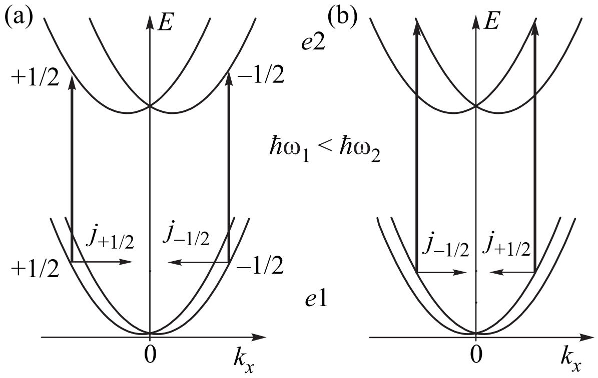


Fig.2. Microscopic origin of pure spin current induced by intersubband photoexcitation. Change of the light frequency (a), (b) leads to reversal of the spin current direction

Due to $\mathbf{k}$-linear spin splitting of the $e1$ and $e2$ subbands, the optical transitions induced by photons of the certain energy $\hbar\omega$ occur only at a fixed $k_x$ where the photon energy matches the energy separation between the subbands [4]. Similarly to the interband excitation considered in the previous section, these $k_x$-points are of opposite sings for the spin branches $\pm 1/2$. Such spin-dependent asymmetry of photoexcitation gives rise to pure spin currents in both $e1$ and $e2$ subbands.

An interesting feature of the pure spin photocurrent induced under intersubband transitions is its spectral behavior: an increase in the photon energy $\hbar\omega$ (see Figs.2a and 2b) leads to a shift of the points $k_x$ that results in reversal of the spin current direction. The explicit spectral dependence of the spin photocurrent in ideal QWs depends on the specific fine structure of the energy spectrum. However, in real structures the spectral width of the intersubband resonance is broadened and hence substantially exceeds the spectral width of the absorption spectrum of an ideal structure. The broadening can be taken into account assuming, e.g., that the energy separation $E_{21}$ between the subbands varies in the QW plane. Then to the first order in the spin-orbit coupling the pure spin current generated under intersubband optical transitions has the form

$$F_{\beta}^{\alpha} = \frac{1}{2\hbar}\left(\gamma_{\alpha\beta}^{(e2)} - \gamma_{\alpha\beta}^{(e1)}\right)\frac{I}{\hbar\omega} \times \\ \times \left[\tau_{e2}\,\eta_{21}(\hbar\omega) + (\tau_{e1}-\tau_{e2})\,\bar{E}\,\frac{d\eta_{21}(\hbar\omega)}{d\,\hbar\omega}\right], \qquad (5)$$

where $\gamma_{\alpha\beta}^{(e1)}$ and $\gamma_{\alpha\beta}^{(e2)}$ are the constants of the spin-orbit coupling in the $e1$ and $e2$ subbands, $\tau_{e1}$ and $\tau_{e2}$ are the corresponding relaxation times of the spin currents, $\eta_{21}(\hbar\omega)$ is the intersubband light absorbance with the inhomogeneous broadening being taken into account, and $\bar{E}$ is the mean value of the electron kinetic energy. It equals to $E_F/2$ for a 2D degenerate gas with the Fermi energy $E_F$ and $k_B T$ for a 2D non-degenerate gas at the temperature $T$. Intersubband light absorption is dominated by spin-conserving optical transitions, therefore the pure spin current (5) is proportional to the difference of subband spin splittings. The spectral behavior of the spin current is determined mainly by the second term in Eq. (5) and repeats the derivative of the light absorption spectrum $d\eta_{21}(\hbar\omega)/\hbar\omega$, since the relaxation time in the excited subband $\tau_{e2}$ is usually shorter than that in the lowest subband, $\tau_{e1}$.

**4. Free-carrier absorption.** Light absorption by free carriers, or the Drude-like absorption, occurs in doped semiconductor structures when the photon energy $\hbar\omega$ is smaller than the band gap as well as the intersubband spacing. Because of the energy and momentum conservation the free-carrier optical transitions become possible if they are accompanied by electron scattering by acoustic or optical phonons, static defects etc. Scattering-assisted photoexcitation with linearly- or unpolarized light also gives rise to a pure spin current. However, in contrast to the direct transitions considered above, the spin splitting of the energy spectrum leads to no essential contribution to the spin current induced by free-carrier absorption. The more important contribution comes from asymmetry of the electron spin-conserving scattering. In semiconductor QWs the matrix element $V$ of electron scattering by static defects or phonons has, in addition to the main contribution $V_0$, an asymmetric spin-dependent term [14−17]

$$V = V_0 + \sum_{\alpha\beta} V_{\alpha\beta}\,\sigma_{\alpha}(k_{\beta} + k'_{\beta})\,, \qquad (6)$$

where $\mathbf{k}$ and $\mathbf{k}'$ are the electron initial and scattered wave vectors, respectively. Microscopically this contribution is caused by the structural and bulk inversion asymmetry similar to the Rashba/Dresselhaus spin splitting of the electron subbands. The asymmetry of the electron-phonon interaction results in non-equal rates of indirect optical transitions for opposite wave vectors in each spin branch. This is illustrated in Fig.3, where

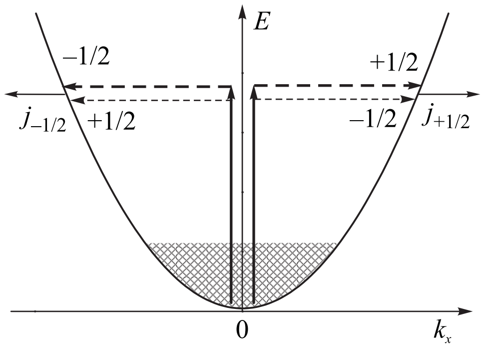


Fig.3. Microscopic origin of pure spin current induced under light absorption by free electrons. The free-carrier absorption is a combined process involving electron-photon interaction (vertical solid lines) and electron scattering (dashed horizontal lines)

the free-carrier absorption is shown as a combined two-stage process involving electron-photon interaction (vertical solid lines) and electron scattering (dashed horizontal lines). The scattering asymmetry is shown by thick and thin dashed lines: electrons with the spin $+1/2$ are preferably scattered into the states with $k_x > 0$, while particles with the spin $-1/2$ are scattered predominantly into the states with $k_x < 0$. The asymmetry causes an imbalance in the distribution of photoexcited carriers in each branch $s = \pm 1/2$ over the positive and negative $k_x$ states and yields oppositely directed electron flows $\mathbf{j}_{\pm 1/2}$ shown by horizontal arrows. Similarly to the interband excitation considered in the previous section, this non-equilibrium distribution is characterized by a pure spin current without charge transfer.

If the photon energy $\hbar\omega$ exceeds the typical electron kinetic energy $\bar{E}$, the pure spin current induced by free-carrier light absorption is given by

$$F_x^\alpha = \frac{\tau_e}{\hbar}\left[\frac{V_{\alpha x}}{V_0}\left(1+\frac{|e_x|^2-|e_y|^2}{2}\right)+\frac{V_{\alpha y}}{V_0}\,e_x e_y\right]\eta_{e1} I\ , \tag{7}$$

where $\eta_{e1}$ is the light absorbance in this spectral range. The components $F_y^\alpha$ are obtained from Eq. (7) by replacement $x \leftrightarrow y$.

In addition to the free-carrier absorption, the spin-dependent asymmetry of electron-phonon interaction can also give rise to a pure spin current in the process of energy relaxation of the photoexcited electrons. In this relaxational mechanism the spin current is generated in a system of hot carriers, independently of heating means.

Mechanisms of pure spin photocurrent considered above can reveal itself in an appearance of electric current in the presence of an in-plane magnetic field. Indeed, the application of a magnetic field results, due to the Zeeman effect, in different equilibrium populations of the subbands. The currents $\mathbf{j}_{\pm 1/2}$ flowing in the opposite directions become non-equivalent which results in a spin polarized net electric current [18].

*Valley-orbit current.* In addition to the spin, free charge carriers can be characterized by another internal property, e.g., by a well number in multi-QW structures or a valley index, $l$, in multi-valley semiconductors. Thus, one can consider not only pure spin currents but also pure orbit-valley currents in which case the net electric current $\mathbf{j} = \sum_l \mathbf{j}_l$ vanishes but the partial currents $\mathbf{j}_l$ contributed by carriers in the $l$th valley are nonzero. In carbon nanotubes the index $l$ runs through two equivalent one-dimensional subbands $(n, K)$ and $(-n, K')$ formed near the $K$ and $K'$ valleys of the graphene sheet rolled-up into a cylinder, where $n$ is the component of orbital angular momentum along the tube principal axis $z$ [19]. In chiral nanotubes the photoexcitation results in nonzero partial flows, $j_z(K)$ and $j_z(K')$, which have opposite signs for linearly polarized light. Another example is a GeSi/Si (111)-grown quantum well structure. It has the overall $C_{3v}$ symmetry and contains three equivalent two-dimensional valleys $l = 1, 2, 3$. The symmetry representing an individual valley is reduced to $C_s$ and allows generation of a partial in-plane photocurrent $\mathbf{j}_l$ under normal light incidence. The net electric current is absent but one can introduce the pure valley-orbit flows $\mathbf{j} = \sum_l c_l \mathbf{j}_l$, where $c_l$ are arbitrary nonequal coefficients.

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