Spin-dependent resonant tunneling in symmetrical double-barrier structure

M.M. Glazov, P.S. Alekseev, M.A. Odoblyudov, V.M. Chistyakov, S.A. Tarasenko and I.N. Yassievich
A.F. Ioffe Physico-Technical Institute, RAS, 194021 St. Petersburg, Russia
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A theory of resonant spin-dependent tunneling has been developed for symmetrical double-barrier structures grown of non-centrosymmetrical semiconductors. The dependence of the tunneling transparency on the spin orientation and the wave vector of electrons leads to (i) spin polarization of the transmitted carriers in an in-plane electric field, (ii) generation of an in-plane electric current under tunneling of spin-polarized carriers. These effects originated from spin-orbit coupling-induced splitting of the resonant level have been considered for double-barrier tunneling structures.

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I. INTRODUCTION

Physics of spin-dependent tunneling phenomena in semiconductor structures has attracted lately a great deal of attention. Significant progress has been made in experimental and theoretical investigation of spin-polarized transport in magnetic tunneling junctions (for review see [1]). On the other hand, it was pointed out recently that electron tunneling could be spin-dependent even in the case of nonmagnetic barriers. It was demonstrated that the transparency of a semiconductor barrier depends on the spin orientation of carriers if the system lacks a center of inversion. Two microscopic mechanisms, Rashba spin-orbit coupling induced by the barrier asymmetry and the Dresselhaus spin-orbit splitting in non-centrosymmetrical materials were shown to be responsible for the effect of spin-dependent tunneling. Spin-orbit interaction couples spin states and space motion of conduction electrons that opens a possibility to orient, manipulate and detect spins by electrical means. Effect of spin-dependent tunneling in nonmagnetic semiconductor heterostructures was supposed to be applied for spin injection and detection of spin-polarized carriers. Devices based on spin-dependent tunneling were suggested to be utilized as components of the spin field-effect transistor.

In this paper we present a theory of spin-dependent tunneling through a double-barrier structure, the structure of resonant tunneling diode (RTD), grown of non-centrosymmetrical semiconductors. Section II is devoted to calculation of the spin-dependent transmission coefficient. The dependence of the structure transparency on the spin orientation and the wave vector of electrons can be employed for spin injection and detection: (i) an electric current flow in the plane of interfaces leads to the spin polarization of the transmitted carriers, (ii) transmission of the spin-polarized carriers is accompanied by generation of an in-plane electric current. These effects are considered in Sections III and IV, respectively. The results of the numerical calculations are compared with that obtained in a simple analytical theory.

II. SPIN-DEPENDENT RESONANT TUNNELING

We consider the transmission of electrons with the initial wave vector $k = (k_x, k_z)$ through a symmetrical double barrier structure grown along $z || [001]$ direction (see Fig. 1). Here $k_x$ is the wave vector in the plane of the interfaces, and $k_z$ is wave vector component normal to the barrier and pointing in the direction of tunneling. The electron motion in each layer of the structure is described by the Hamiltonian

$$\hat{H} = -\frac{\hbar^2}{2m^*} \frac{\partial^2}{\partial z^2} + \frac{\hbar^2 k_z^2}{2m^*} + V(z) + \hat{H}_D,$$

(1)

where $m^*$ is the electron effective mass, $V(z)$ is the heterostructure potential equal to $V_b$ for the barriers and $-V_w$ for the well, and $\hat{H}_D$ is the spin-dependent $k^3$ Dresselhaus term that describes spin-orbit splitting of the conduction band in zinc-blende-lattice semiconductors. We assume the RTD structure to be designed so that the resonant transmission occurs for the incident electrons with the kinetic energy $\varepsilon$ much smaller than $V_b$ and $V_w$. Then the Dresselhaus term is simplified to

$$\hat{H}_D = \gamma (\hat{\sigma}_x k_x - \hat{\sigma}_y k_y) \frac{\partial^2}{\partial z^2},$$

(2)

where $\gamma$ is a material constant, $\hat{\sigma}_x$ and $\hat{\sigma}_y$ are the Pauli matrices, and the coordinate axes $x, y, z$ are assumed to

![FIG. 1: Electron transmission through (001)-grown RTD structure. $V_b$ is the height of the barriers, $V_w$ is the depth of the well, $a$ and $b$ are the well width and barrier thickness, respectively.](image-url)
be parallel to the cubic crystallographic axes [100], [010], [001], respectively.

The Dresselhaus term \( \hat{H}_D \) is diagonalized by the spinors

\[
\chi_{\pm} = \frac{1}{\sqrt{2}} \left( 1 \pm 1 \text{e}^{-i\varphi} \right),
\]

which describe the electron states “+” and “−” of the opposite spin directions. Here \( \varphi \) is the polar angle of the wave vector \( k \) in the \( xy \) plane, \( k_{||} = (k_{||} \cos \varphi, k_{||} \sin \varphi) \). The electron spins \( s_{\pm} \) corresponding to the eigen-states “+” are given by

\[
s_{\pm}(k_{||}) = \frac{1}{2} \chi_{\pm}^* \sigma \chi_{\pm} = \frac{1}{2} \left( \mp \cos \varphi, \pm \sin \varphi, 0 \right).
\]

One can note that essentially \( \hat{H}_D \) induces a spin-dependent correction to the effective mass for electron motion along \( z \). In the basis of the spin eigen-states “+” the effective Hamiltonian \( \hat{H}_D \) has the simple form

\[
H_{\pm} = -\frac{\hbar^2}{2m_{\pm}} \partial^2 + \frac{\hbar^2 k_{||}^2}{2m_{\pm}} + V(z),
\]

where the effective mass along \( z \), modified by spin-orbit coupling, depends on the in-plane electron wave vector and is given by

\[
m_{\pm} = m^* \left( 1 \pm 2 \frac{\gamma m^* k_{||}^2}{\hbar^2} \right)^{-1}.
\]

Solution of the Schrödinger equation with the Hamiltonian \( \hat{H}_D \) and boundary conditions for the wave functions \( \psi_{\pm} \), which require the continuity of

\[
\psi_{\pm} \quad \text{and} \quad \frac{1}{m_{\pm}} \frac{\partial \psi_{\pm}}{\partial z}
\]

at the interfaces, allows one to derive coefficients of transmission, \( t_{\pm} \), reflection, \( r_{\pm} \), for the electrons of spin eigen-states “+” and “−”.

Fig. 2 presents the dependences of the double-barrier structure transparency \( |t_{\pm}|^2 \) on the incident electron energy along growth direction, \( \varepsilon_z = \hbar^2 k_{||}^2 / 2m^* \), calculated numerically for the fixed in-plane wave vector \( k_{||} \). The spin splitting of the resonant peak is clearly seen. In calculations both the electron effective mass \( m^* \) and the Dresselhaus constant \( \gamma \) are assumed to be the same for the barrier and well layers.

In the limit of thick barriers the structure transparency demonstrates sharp peaks and hence can be approximated by Dirac \( \delta \)-functions

\[
|t_{\pm}(\varepsilon_z, k_{||})|^2 \approx \pi \Gamma_{\pm}(k_{||}) \delta \left[ \varepsilon_z - E_{\pm}(k_{||}) \right],
\]

where the prefactors \( \Gamma_{\pm}(k_{||}) \) describe the transmission efficiencies, and \( E_{\pm}(k_{||}) \) stand for the energies of the resonances. The positions of the resonances \( E_{\pm}(k_{||}) \) correspond to the energies of size-quantization of an electron in the quantum well of infinitely thick barriers with the

\[
\varepsilon_{\pm}(k_{||}) = \frac{\hbar^2 k_{||}^2}{2m^*} + V_0 - \frac{\pi}{2} \Gamma_{\pm}(k_{||}) \delta \left[ \varepsilon_z - E_{\pm}(k_{||}) \right],
\]

Dresselhaus spin-orbit interaction included. The transmission efficiencies \( \Gamma_{\pm}(k_{||}) \) are determined, to the contrary, by the electron lifetimes on the resonant levels in the double-barrier structure. Considering the spin-orbit interaction to be small perturbation, the positions of the resonant levels and their widths can be expanded as

\[
E_{\pm}(k_{||}) = E_0 \pm \alpha k_{||}, \quad \Gamma_{\pm}(k_{||}) = (1 \pm \beta k_{||}) \Gamma_0,
\]

where \( E_0 \) and \( \Gamma_0 \) are the level position and the width when the spin-orbit interaction is neglected,

\[
\Gamma_0 = 8 \sqrt{E_0(V_b - E_0)} \frac{(V_b + E_0)(V_w + E_0)}{V_b(V_b + V_w)} \exp(-2\kappa b) / (1 + \kappa a)^{1/2},
\]

\[
\kappa = \sqrt{2m^*(V_b - E_0)/\hbar} \quad \text{is the reciprocal length of the wave function decay under the barrier,} \quad a \quad \text{is the quantum well width,} \quad b \quad \text{is the thickness of the barriers,} \quad \alpha \quad \text{and} \quad \beta \quad \text{are the coefficients which describe the spin splitting of the level and the width modification of the spin sublevels, respectively. For the case under study,} \quad E_0 \ll V_b, V_w, \quad \text{they are given by}
\]

\[
\alpha = \frac{2 \gamma m^*}{\hbar^2} \frac{V_w}{1 + 2/\kappa a},
\]

\[
\beta = \alpha \left( \frac{k b}{V_b} + \frac{1}{2E_0} \right) + 2 \gamma m^* k b / \hbar^2.
\]

Note, that the width \( \Gamma_0 \) and the coefficient \( \beta \) are highly sensitive to the position of the resonant level \( E_0 \). For the RTD structure parameters presented in the caption to Fig. 2 the coefficients can be estimated as follows: \( \alpha = 9.7 \times 10^{-7} \text{meV} \cdot \text{cm} \), \( \beta = (4.2 \times 10^{-8}+4.8 \times 10^{-7} \cdot \text{meV} / E_0) \text{ cm} \).
III. SPIN ORIENTATION OF CARRIERS

Spin splitting of the resonant peak at non-zero \( k_{\parallel} \) can be employed for injection of spin-polarized carriers. We assume two parts of bulk semiconductor separated by the RTD structure, and electrons tunneling through the barrier from the left to the right side of the structure. In equilibrium the momentum distribution of the incident electrons is isotropic in the interface plane and therefore the average spin of the transmitted carriers vanishes. This isotropy can be broken by application, for example, of an in-plane electric field, \( \mathbf{F} \). Then the carriers tunnel through the structure with non-zero average wave vector in the plane of interfaces that leads to the spin polarization of the transmitted electrons.\(^8\)

The average spin of the transmitted electrons is given by

\[
\mathbf{s} = \mathbf{\dot{S}}/\mathbf{\dot{N}},
\]

where \( \mathbf{\dot{S}} \) and \( \mathbf{\dot{N}} \) are the spin and carrier fluxes through the barrier given in the linear in \( \mathbf{F} \) regime by

\[
\mathbf{\dot{S}} = \sum_{k_{\perp}, k_{z} > 0} f_{1}(\mathbf{k}) \left[ |t_{+}(\varepsilon_{z}, k_{\parallel})|^{2} s_{+}(k_{\parallel}) + |t_{-}(\varepsilon_{z}, k_{\parallel})|^{2} s_{-}(k_{\parallel}) \right] v_{z},
\]

\[
\mathbf{\dot{N}} = \sum_{k_{\perp}, k_{z} > 0} f_{0}(\varepsilon) \left[ |t_{+}(\varepsilon_{z}, k_{\parallel})|^{2} s_{+}(k_{\parallel}) |t_{-}(\varepsilon_{z}, k_{\parallel})|^{2} s_{-}(k_{\parallel}) \right] v_{z},
\]

\( \mathbf{v} = \hbar \mathbf{k}/m^{*} \) is the electron velocity, \( f_{0}(\varepsilon) \) is the equilibrium distribution function, \( \varepsilon \) is the electron energy, \( f_{1}(\mathbf{k}) \) is the electric field-induced correction to the distribution function,

\[
f_{1}(\mathbf{k}) = -e\tau_{p} \frac{d f_{0}}{d\varepsilon} (\mathbf{v}_{d} \cdot \mathbf{F}),
\]

\( e \) is the electron charge, and \( \tau_{p} \) is the momentum relaxation time.

Substituting the spin vectors of the Dresselhaus eigenstates \( \mathbf{s}_{\pm} \) in the form Eq. (11) into Eq. (12), one derives the angular dependence of the average spin of the transmitted carriers,

\[
s_{x} = \frac{v_{d,x}}{2 v_{d}} P_{s}, \quad s_{y} = -\frac{v_{d,y}}{2 v_{d}} P_{s},
\]

where \( v_{d} = (e\tau_{p}/m^{*})\mathbf{F} \) is the in-plane drift velocity of the incident electrons, and \( P_{s} \) is the spin polarization of the transmitted particles,

\[
P_{s} = \frac{v_{d} m^{*}}{2 N} \sum_{k_{\perp}, k_{z} > 0} \frac{d f_{0}}{d\varepsilon} \left[ |t_{+}(\varepsilon_{z}, k_{\parallel})|^{2} - |t_{-}(\varepsilon_{z}, k_{\parallel})|^{2} \right] v_{z} v_{||}.
\]

The direction of the electron spin is determined by the symmetry of the Dresselhaus term. In particular, the spin \( \mathbf{s} \) is parallel (or antiparallel) to the electron drift velocity \( \mathbf{v}_{d} \), if \( \mathbf{v}_{d} \) is directed along the crystal cubic axis \([100]\) or \([010]\); and \( \mathbf{s} \) is perpendicular to \( \mathbf{v}_{d} \), if the latter is directed along the the axis \([110]\) or \([101]\).

In the limit of thick barriers the structure transparency can be approximated by Dirac \( \delta \)-functions. Then substituting the transparency \( |t_{\parallel}|^{2} \) in the form Eq. (11) and assuming spin-orbit interaction to be small, one derives the following expression for the spin polarization of the transmitted carriers

\[
P_{s} = \frac{v_{d} m^{*}}{2 \hbar} \left( \alpha/\zeta - \beta \right),
\]

where \( \zeta = \int_{E_{0}}^{\infty} f_{0}(\varepsilon) d\varepsilon / f_{0}(E_{0}) \) is an energy equal to \( E_{F} - E_{0} \) for 3D Fermi and \( k_{B}T \) for 3D Boltzmann gas, respectively.

The dependence of the spin polarization \( P_{s} \) on the resonant level energy \( E_{0} \) is presented in Fig. 3 for the incident electrons forming degenerate gas with the Fermi energy \( E_{F} = 10 \text{ meV} \) and the gas of the same carrier concentration at temperature \( T = 300 \text{ K} \). Solid curves correspond to the exact solution of Eq. (14), with the transmission coefficients \( t_{\pm}(\varepsilon, k_{\parallel}) \) being calculated numerically in the framework of the transfer matrices. The variation of \( E_{0} \) from 0 to 10 meV was achieved modifying the quantum well width from \( a \approx 31 \text{ Å} \) to \( a \approx 29.4 \text{ Å} \). Dashed curves are plotted following Eq. (15) for Fermi and Boltzmann statistics. One can see that in the wide range of the resonant level positions the simple analytical theory demonstrates a good agreement with the numerical calculations. For the reasonable set of parameters given in the figure caption it is possible to achieve spin polarizations of several percents. The sign of the po-

![FIG. 3: The spin polarization \( P_{s} \) as a function of the resonant level position \( E_{0} \) for degenerate electron gas with \( E_{F} = 10 \text{ meV} \) and non-degenerate gas of the same carrier concentration at \( T = 300 \text{ K} \). Solid curves correspond to the numerical calculation, dashed curves are plotted following Eq. (15). The parameters of the double-barrier structure are presented in the caption to Fig. 2 and \( v_{d} = 2.5 \times 10^{6} \text{ cm/s} \) that corresponds to \( v_{F}/10 \).](image-url)
larization $P_e$ is governed by interplay between the spin splitting, $\alpha k_{\parallel}$, and the difference of the widths of the spin sublevels, $\beta k_{\parallel}$, since the carrier occupation of the lower sublevel "−" is larger than that of the higher sublevel "+" while the tunneling transparency of the sublevel "−" is smaller than that of "+" (see Eq. (7)). It clarifies why the terms proportional to $\alpha$ and $\beta$ contribute to Eq. (15) with opposite signs. At low temperature both terms are comparable, and $P_e$ changes the sign with increasing of $E_0$. At $E_0 \ll E_F$ the effect is mainly determined by the difference of the spin sublevel transparencies. The role of the energy spacing between the sublevels is negligible, since the carrier populations at the energies $E_+$ and $E_−$ almost coincide. With increasing of $E_0$ the "step-like" Fermi distribution leads to different occupations of the spin sublevels, and the polarization $P_e$ changes the sign. At higher temperatures the carrier distribution becomes to be smooth, and the effect is related mainly to the difference of the tunneling transparencies of the spin sublevels.

IV. TUNNELING SPIN GALVANIC EFFECT

Generation of an electric current by spin-polarized carriers represents the effect inverse to spin injection. Now we assume the electron gas on the left side of the double-barrier structure to be spin-polarized. Electrons with various wave vectors tunnel through the RTD. However, due to the splitting of the resonant level the flux of the spin-polarized carriers on the left side of the structure to form 3D non-degenerate gas at the temperature $T$. The spin matrix of the electron transmission through the structure is given by

$$T = t_+(\varepsilon, k_{||}) \chi_+ \chi_+^\dagger + t_-(\varepsilon, k_{||}) \chi_- \chi_-^\dagger . \quad (18)$$

Substituting the density matrix (16) and the transmission matrix (15) into Eq. (10) and taking into account the definition of the vectors $s_{\pm}$ (1), one obtains

$$j_\parallel = \frac{4e \tau_p \rho_s}{\langle 1/\varepsilon \rangle} \sum_{k_{||}, k_z > 0} \frac{df_0}{d\varepsilon} \left[ t_+(\varepsilon, k_{||}) \right]^2 n_s \cdot s_+(k_{||})$$

$$+ \left| t_-(\varepsilon, k_{||}) \right|^2 n_s \cdot s_-(k_{||}) \right] v_\parallel v_z . \quad (19)$$

Taking into account the explicit form of the spin vectors of the Dresselhaus eigen-states (1), the components of the tunneling spin-galvanic current have the form

$$j_{||, x} = -j_{||, y} = j_{||, z} = j_{||} n_s, \quad (20)$$

where

$$j_{||} = -\frac{e \tau_p \rho_s}{\langle 1/\varepsilon \rangle} \sum_{k_{||}, k_z > 0} \frac{df_0}{d\varepsilon} \left[ |t_+(\varepsilon, k_{||})|^2 - |t_-(\varepsilon, k_{||})|^2 \right] v_z v_z . \quad (21)$$

The direction of the tunneling spin-galvanic current is determined by the spin orientation of the electrons with respect to the crystallographic axes.

In the limit of thick barriers the structure transparency can be obtained substituting $\delta$-functions (4) for the structure transparency. Then assuming spin-orbit interaction to be small, the interface current is derived to be

$$j_{||} = -\frac{e \tau_p \rho_s}{\langle 1/\varepsilon \rangle} \frac{f(E_0) m^*}{\pi \hbar^2} (\alpha - \zeta \beta) \Gamma_0 . \quad (22)$$

Fig. 1 presents the dependence of the interface current $j_{||}$ on the energy position of the resonant level $E_0$ for the incident electrons forming spin-polarized degenerate gas with the Fermi energy $E_F = 10$ meV and the gas of the same carrier concentration at temperature $T = 300$ K. Solid curves correspond to the exact solution of Eq. (21), with the transmission coefficients $t_\pm(\varepsilon, k_{||})$ being calculated numerically. Dashed curves are plotted following Eq. (22) for Fermi and Boltzmann statistics. For the AlGaSb-based RTD structure considered as an example the tunneling spin-galvanic current is of order of mA/cm. Estimation shows that it is enhanced by an order of magnitude with respect to the interface current generated under tunneling through a single AlGaSb-based barrier provided the equal electron tunneling flux,
FIG. 4: The interface current $j_\parallel$ as a function of the resonant level position $E_0$ for degenerate electron gas with $E_F = 10$ meV and non-degenerate gas of the same carrier concentration at $T = 300$ K. Solid curves correspond to the numerical calculation, dashed curves are plotted following Eq. (22). The parameters of the double-barrier structure are presented in the caption to Fig 2, $\tau = 1$ ps, and $p_s = 0.1$.

In conclusion, the theory of spin-dependent electron tunneling has been developed for symmetrical double-barrier structures based on zinc-blende-lattice semiconductor compounds. The Dresselhaus spin-orbit interaction couples spin states and space motion of conduction electrons that leads to spin splitting of the resonant level depending on the in-plane electron wave vector. The effect of the spin-dependent tunneling could be employed for creating spin injectors and detectors based on non-magnetic tunneling structures.

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* Electronic address: tarasenko@coherent.ioffe.ru

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