Spin-dependent resonant tunneling in symmetrical double-barrier structures

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A theory of resonant spin-dependent tunneling has been developed for symmetrical double-barrier structures grown of noncentrosymmetrical 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 and (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 a great deal of attention lately. Significant progress has been made in experimental and theoretical investigation of spin-polarized transport in magnetic tunneling junctions (for review see Refs. 1 and 2). 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[3–5] and the k^3 Dresselhaus spin-orbit splitting in noncentrosymmetrical materials,[6,7] 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[8–11] and detection of spin-polarized carriers.[12,13] Devices based on spin-dependent tunneling were suggested to be utilized as components of the spin field-effect transistor.[12]

In this paper we present a theory of spin-dependent tunneling through a double-barrier structure, the structure of a resonant tunnel diode (RTD), grown of noncentrosymmetrical 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 and (ii) transmission of the spin-polarized carriers is accompanied by generation of an in-plane electric current. These effects are considered in Secs. 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 $\boldsymbol{k}=(\boldsymbol{k}_\parallel,k_z)$ through a symmetrical double barrier structure grown along the $z\parallel[001]$ direction (see Fig. 1). Here $\boldsymbol{k}_\parallel$ 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 tunnelling. The electron motion in each layer of the structure is described by the Hamiltonian

$$\hat{H}=-\frac{\hbar^2}{2m^*}\frac{\partial^2}{\partial z^2}+\frac{\hbar^2k_\parallel^2}{2m^*}+V(z)+\hat{\mathcal{H}}_D, \tag{1}$$

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 $\hat{\mathcal{H}}_D$ 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 ε much smaller than V_b and V_w. Then the Dresselhaus term is simplified to[7]

$$\hat{\mathcal{H}}_D=\gamma(\hat{\sigma}_xk_x-\hat{\sigma}_yk_y)\frac{\partial^2}{\partial z^2}, \tag{2}$$

where γ 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 be parallel to the cubic crystallographic axes [100], [010], [001], respectively.

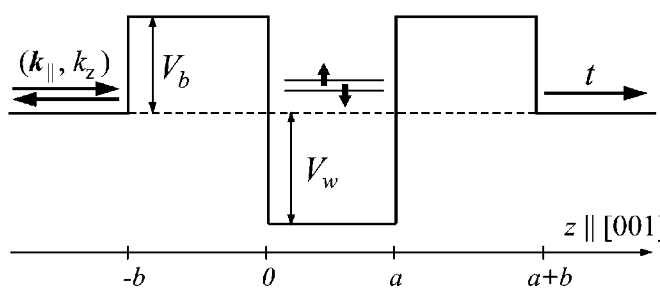

FIG. 1. Electron transmission through (001)-grown RTD structure. V_b is the height of the barriers, V_w is the depth of the well, a and b are the well width and barrier thickness, respectively.

The Dresselhaus term (2) is diagonalized by the spinors

$$\chi_\pm = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 \\ \mp e^{-i\varphi} \end{pmatrix}, \qquad (3)$$

which describe the electron states "+" and "−" of the opposite spin directions. Here φ is the polar angle of the wave vector $\boldsymbol{k}$ in the xy plane, $\boldsymbol{k}_\parallel = (k_\parallel \cos\varphi, k_\parallel \sin\varphi)$. The electron spins $\boldsymbol{s}_\pm$ corresponding to the eigenstates "±" are given by

$$\boldsymbol{s}_\pm(\boldsymbol{k}_\parallel) = \frac{1}{2}\chi_\pm^\dagger \hat{\boldsymbol{\sigma}} \chi_\pm = \frac{1}{2}(\mp\cos\varphi, \pm\sin\varphi, 0). \qquad (4)$$

One can note that $\hat{\mathcal{H}}_D$ essentially induces a spin-dependent correction to the effective mass for electron motion along z. In the basis of the spin eigenstates "±" the effective Hamiltonian (1) has the simple form

$$H_\pm = -\frac{\hbar^2}{2m_\pm}\frac{\partial^2}{\partial z^2} + \frac{\hbar^2 k_\parallel^2}{2m^*} + V(z), \qquad (5)$$

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

$$m_\pm = m^*\left(1 \pm 2\frac{\gamma m^* k_\parallel}{\hbar^2}\right)^{-1}. \qquad (6)$$

Solution of the Schrödinger equation with the Hamiltonian (5) and boundary conditions for the wave functions $\psi_\pm$, which require the continuity of

$$\psi_\pm \text{ and } \frac{1}{m_\pm}\frac{\partial\psi_\pm}{\partial z}$$

at the interfaces, allows one to derive coefficients of transmission $t_\pm$ and reflection $r_\pm$ for the electrons of spin eigenstates "+" and "−."

Figure 2 presents the dependencies of the double-barrier structure transparency $|t_\pm|^2$ on the incident electron energy along growth direction $\varepsilon_z = \hbar^2 k_z^2/2m^*$, calculated numerically for the fixed in-plane wave vector $k_\parallel$. The spin splitting of the resonant peak is clearly seen. In calculations both the electron effective mass m^* and the Dresselhaus constant γ 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 δ functions

$$|t_\pm(\varepsilon_z, k_\parallel)|^2 \approx \pi\Gamma_\pm(k_\parallel)\delta[\varepsilon_z - E_\pm(k_\parallel)], \qquad (7)$$

where the prefactors $\Gamma_\pm(k_\parallel)$ describe the transmission efficiencies and $E_\pm(k_\parallel)$ stand for the energies of the resonances. The positions of the resonances $E_\pm(k_\parallel)$ correspond to the energies of size quantization of an electron in the quantum well of infinitely thick barriers with the Dresselhaus spin-orbit interaction included. The transmission efficiencies $\Gamma_\pm(k_\parallel)$ are determined, on 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

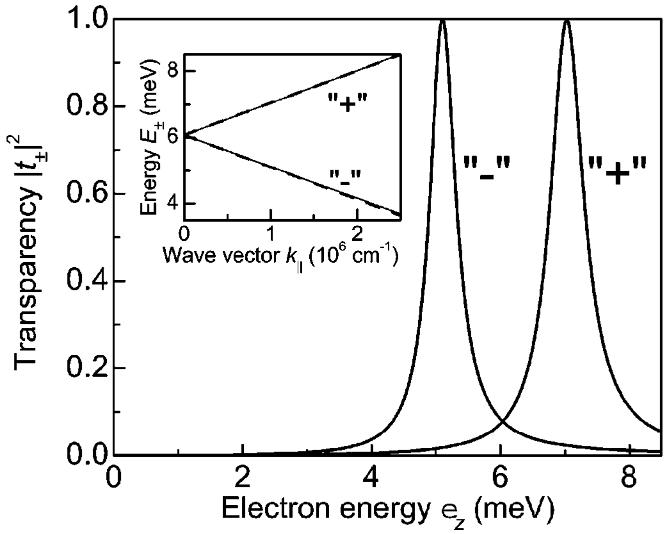

FIG. 2. The transparency of the double-barrier structure $|t_\pm|^2$ as a function of ε_z at fixed $k_\parallel = 10^6$ cm^{-1}. The insert shows the dependence of the spin splitting of the resonant peak on $k_\parallel$ calculated numerically (solid curves) and following Eq. (8) (dashed lines). The used parameters, $\gamma = 76$ eV Å^3, $m^* = 0.053 m_0$, $V_b = 230$ meV, $V_w = 200$ meV, $a = 30$ Å, and $b = 50$ Å, correspond to $Al_xGa_{1-x}Sb$, $x = 0.15/0.3/0/0.3/0.15$, RTD structure. (Refs. 14–16).

$$E_\pm(k_\parallel) = E_0 \pm \alpha k_\parallel, \quad \Gamma_\pm(k_\parallel) = (1 \pm \beta k_\parallel)\Gamma_0, \qquad (8)$$

where E_0 and Γ_0 are the level position and the width when the spin-orbit interaction is neglected,

$$\Gamma_0 = 8\sqrt{E_0(V_b - E_0)}\frac{(V_b - E_0)(V_w + E_0)}{V_b(V_b + V_w)}\frac{\exp(-2\kappa b)}{1 + \kappa a/2}, \qquad (9)$$

$\kappa = \sqrt{2m^*(V_b - E_0)}/\hbar$ is the reciprocal length of the wave function decay under the barrier, a is the quantum well width, b is the thickness of the barriers, and α and β 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, $E_0 \ll V_b, V_w$, they are given by

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

$$\beta = \alpha\left(\frac{\kappa b}{V_b} + \frac{1}{2E_0}\right) + 2\gamma\frac{m^*\kappa b}{\hbar^2}. \qquad (10)$$

Note, that the width Γ_0 and the coefficient β 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: $\alpha = 9.7 \times 10^{-7}$ meV cm, $\beta = (4.2 \times 10^{-8} + 4.8 \times 10^{-7}$ meV$/E_0)$ cm.

III. SPIN ORIENTATION OF CARRIERS

Spin splitting of the resonant peak at nonzero $\boldsymbol{k}_\parallel$ 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 $\boldsymbol{F}$. Then the carriers tunnel through the structure with nonzero average wave vector in the plane of interfaces that leads to the spin polarization of the transmitted electrons.[8,9]

The average spin of the transmitted electrons is given by

$$\boldsymbol{s} = \dot{\boldsymbol{S}}/\dot{N}, \tag{11}$$

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

$$\dot{\boldsymbol{S}} = \sum_{\boldsymbol{k}_\parallel, k_z>0} f_1(\boldsymbol{k})[|t_+(\varepsilon_z,k_\parallel)|^2 \boldsymbol{s}_+(\boldsymbol{k}_\parallel) + |t_-(\varepsilon_z,k_\parallel)|^2 \boldsymbol{s}_-(\boldsymbol{k}_\parallel)]v_z,$$

$$\dot{N} = \sum_{\boldsymbol{k}_\parallel, k_z>0} f_0(\varepsilon)[|t_+(\varepsilon_z,k_\parallel)|^2 + |t_-(\varepsilon_z,k_\parallel)|^2]v_z, \tag{12}$$

$\boldsymbol{v}=\hbar\boldsymbol{k}/m^*$ is the electron velocity, $f_0(\varepsilon)$ is the equilibrium distribution function, ε is the electron energy, $f_1(\boldsymbol{k})$ is the electric field-induced correction to the distribution function,

$$f_1(\boldsymbol{k}) = -e\tau_p \frac{df_0}{d\varepsilon}(\boldsymbol{v}_\parallel \cdot \boldsymbol{F}),$$

e is the electron charge, and τ_p is the momentum relaxation time.

Substituting the spin vectors of the Dresselhaus eigenstates $\boldsymbol{s}_\pm$ in the form (4) into Eq. (12), one derives the angular dependence of the average spin of the transmitted carriers

$$s_x = \frac{v_{d,x}}{2v_d}P_s, \quad s_y = -\frac{v_{d,y}}{2v_d}P_s, \tag{13}$$

where $\boldsymbol{v}_d=(e\tau_p/m^*)\boldsymbol{F}$ is the in-plane drift velocity of the incident electrons and P_s is the spin polarization of the transmitted particles

$$P_s = \frac{v_d m^*}{2\dot{N}} \sum_{\boldsymbol{k}_\parallel, k_z>0} \frac{df_0}{d\varepsilon}[|t_+(\varepsilon_z,k_\parallel)|^2 - |t_-(\varepsilon_z,k_\parallel)|^2]v_z v_\parallel. \tag{14}$$

The direction of the electron spin (13) is determined by the symmetry of the Dresselhaus term. In particular, the spin $\boldsymbol{s}$ is parallel (or antiparallel) to the electron drift velocity $\boldsymbol{v}_d$, if $\boldsymbol{v}_d$ is directed along the crystal cubic axis [100] or [010]; and $\boldsymbol{s}$ is perpendicular to $\boldsymbol{v}_d$, if the latter is directed along the axis $[1\bar{1}0]$ or [110].

In the limit of thick barriers the structure transparency can be approximated by Dirac δ functions. Then substituting the transparency $|t_\pm|^2$ in the form Eq. (7) and assuming spin-orbit interaction to be small, one derives the following expression for the spin polarization of the transmitted carriers

$$P_s = \frac{v_d m^*}{\hbar}(\alpha/\zeta - \beta), \tag{15}$$

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_BT for 3D Boltzmann gas, respectively.

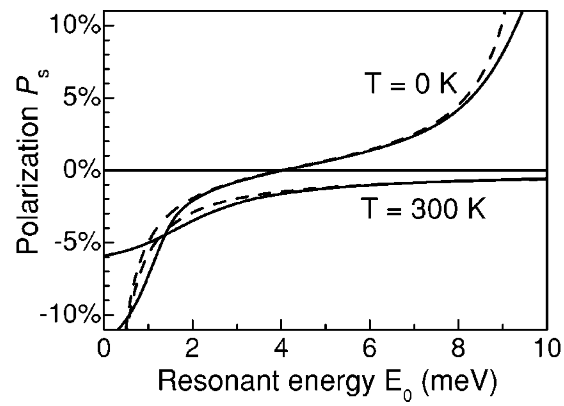

FIG. 3. The spin polarization P_s as a function of the resonant level position E_0 for degenerate electron gas with $E_F=10$ meV and nondegenerate gas of the same carrier concentration at $T=300$ K. Solid curves correspond to the numerical calculation, dashed curves are plotted following Eq. (15). The parameters of the double-barrier structure are presented in the caption to Fig. 2, and $v_d=2.5\times10^6$ cm/s that corresponds to $v_F/10$.

The dependence of the spin polarization P_s on the resonant level position E_0 is presented in Fig. 3 for the incident electrons forming 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 $t_\pm(\varepsilon_z,k_\parallel)$ 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 $a\approx31$ to 29.4 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 polarization P_s 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 α and β contribute to Eq. (15) with opposite signs. At low temperature both terms are comparable, and P_s 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 "steplike" Fermi distribution leads to different occupations of the spin sublevels, and the polarization P_s 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 tunnel flux of the spin-polarized carriers with the certain in-plane wave vector $\boldsymbol{k}_\parallel$ is larger than the flux of the particles with the opposite in-plane wave vector $-\boldsymbol{k}_\parallel$. This asymmetry results in the in-plane flow of the transmitted electrons near the barrier, i.e., in the interface electric current. The direction of this interface current is determined by the spin orientation of the electrons and symmetry properties of the barrier, in particular the current reverses its direction if the spin orientation changes the sign.[12,13]

The theory of such "tunneling spin-galvanic effect" is developed by using the spin density matrix technique. The interface current of spin-polarized electrons transmitted through the tunneling structure is given by[13]

$$\boldsymbol{j}_\parallel = e \sum_{\boldsymbol{k}_\parallel, k_z>0} \tau_p \boldsymbol{v}_\parallel v_z \mathrm{Tr}[\mathcal{T}\rho_l \mathcal{T}^\dagger], \tag{16}$$

where τ_p is the momentum relaxation time, ρ_l is the electron density matrix on the left side of the structure, and $\mathcal{T}$ is the spin matrix of the tunneling transmission that links the incident spinor wave function ψ_l to the transmitted spinor wave function ψ_r, $\psi_r = \mathcal{T}\psi_l$. We assume the carriers on the left side of the structure to form 3D spin-oriented electron gas, and electron distributions in both spin subband to be thermalized. For the case of small degree of spin polarization, the density matrix has the form

$$\rho_l \approx f_0 \hat{I} - \frac{df_0}{d\varepsilon}\frac{2p_s}{\langle 1/\varepsilon\rangle}(\boldsymbol{n}_s \cdot \hat{\sigma}), \tag{17}$$

where f_0 is the equilibrium distribution function of nonpolarized carriers, $\boldsymbol{n}_s$ is the unit vector directed along the spin orientation, p_s is the degree of the polarization, and $\langle 1/\varepsilon\rangle$ is the average value of the reciprocal kinetic energy of the carriers equal to $3/E_F$ for 3D degenerate electron gas with the Fermi energy E_F, and $2/k_B T$ and 3D nondegenerate gas at the temperature T. The spin matrix of the electron transmission through the structure is given by

$$\mathcal{T} = t_+(\varepsilon_z, k_\parallel)\chi_+\chi_+^\dagger + t_-(\varepsilon_z, k_\parallel)\chi_-\chi_-^\dagger. \tag{18}$$

Substituting the density matrix (17) and the transmission matrix (18) into Eq. (16) and taking into account the definition of the vectors $\boldsymbol{s}_\pm$ (4), one obtains

$$\boldsymbol{j}_\parallel = -\frac{4e\tau_p p_s}{\langle 1/\varepsilon\rangle}\sum_{\boldsymbol{k}_\parallel, k_z>0}\frac{df_0}{d\varepsilon}[|t_+(\varepsilon_z,k_\parallel)|^2 \boldsymbol{n}_s\cdot\boldsymbol{s}_+(\boldsymbol{k}_\parallel) + |t_-(\varepsilon_z,k_\parallel)|^2 \boldsymbol{n}_s\cdot\boldsymbol{s}_-(\boldsymbol{k}_\parallel)]\boldsymbol{v}_\parallel v_z. \tag{19}$$

Taking into account the explicit form of the spin vectors of the Dresselhaus eigen-states (4), the components of the tunneling spin-galvanic current have the form

$$j_{\parallel,x} = -j_\parallel n_{s,x}, \quad j_{\parallel,y} = j_\parallel n_{s,y}, \tag{20}$$

where

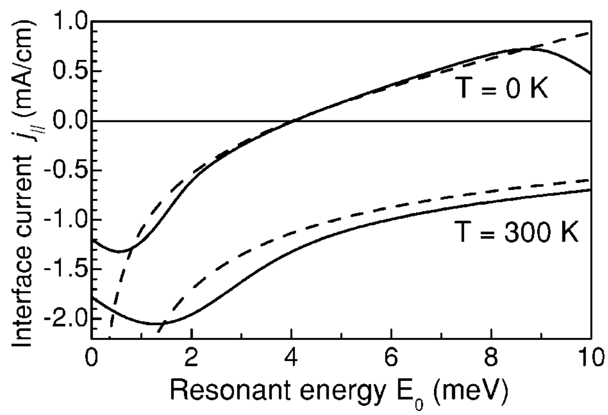

FIG. 4. The interface current $j_\parallel$ as a function of the resonant level position E_0 for degenerate electron gas with $E_F = 10$ meV and nondegenerate gas of the same carrier concentration at $T = 300$ K. Solid curves correspond to the numerical calculation, dashed curves are plotted following Eq. (22). The parameters of the double-barrier structure are presented in the caption to Fig. 2, $\tau_p = 1$ ps and $p_s = 0.1$.

$$j_\parallel = -\frac{e\tau_p p_s}{\langle 1/\varepsilon\rangle}\sum_{k_\parallel, k_z>0}\frac{df_0}{d\varepsilon}[|t_+(\varepsilon_z,k_\parallel)|^2 - |t_-(\varepsilon_z,k_\parallel)|^2]v_z v_\parallel. \tag{21}$$

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 δ functions (7) for the structure transparency. Then assuming spin-orbit interaction to be small, the interface current is derived to be

$$j_\parallel = -\frac{e\tau_p p_s}{\langle 1/\varepsilon\rangle}\frac{f(E_0)m^*}{\pi\hbar^4}(\alpha - \zeta\beta)\Gamma_0. \tag{22}$$

Figure 4 presents the dependence of the interface current $j_\parallel$ 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 $t_\pm(\varepsilon_z, k_\parallel)$ 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 tunnel flux $\dot{N} \sim 10^{22}$ $1/(\mathrm{cm}^2\,\mathrm{s})$.[13] Similarly to the spin polarization (Fig. 3), the sign of the interface current (22) is governed by interplay between the contributions responsible for the spin splitting and the difference of the spin sublevel transparencies.

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 nonmagnetic tunneling structures.

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