Spin-polarized tunneling current between independently contacted quantum wells

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(Received 3 February 2004; published 24 August 2004)

We study the spin-polarized tunneling current between independently contacted quantum wells under an in-plane magnetic field. The splitting of energy spectra of two-dimensional electrons due to both spin-orbit and Pauli interactions is taken into account. The line shape of the resonant peak of the tunneling current is described for both homogeneous and inhomogeneous broadening mechanisms and the effects of temperature and finite drop of voltage is investigated. We show that a considerable spin-polarized current (the degree of polarization about 80%) can be achieved in the InAs-based double-well structures.

DOI: 10.1103/PhysRevB.70.075311 PACS number(s): 72.25.Mk, 73.21.Fg, 73.40.Gk

I. INTRODUCTION

Tunneling current between independently contacted quantum wells (see Refs. 1 and 2 for review) and wires (see references in Ref. 3) has been investigated during the last decade. Due to the restrictions imposed by the energy and momentum conservation laws, the nonzero current between two low-dimensional states with simple parabolic dispersion laws exists only if the electron scattering is taken into account. The tunneling current in this case essentially depends on the scattering rate which determines both the height and width of the resonance tunneling peak. In the presence of a magnetic field parallel to two-dimensional (2D) layers, as shown in Fig. 1(a), the momentum-conserving tunneling, which does not require scattering, takes place[4,5] because of the relative diamagnetic shift leading to intersection of energy spectra of electrons in the wells in the energy-momentum space. The tunneling current in this case is less sensitive to the scattering and reflects essential features of the electron energy spectrum.

A similar momentum-conserving regime of tunneling in the absence of the magnetic field exists in the systems with spin splitting of electron states due to spin-orbit interaction. The cases of tunneling between spin-split one-dimensional (1D) electron states[6] and 2D electron states[7] has been recently considered theoretically. The tunneling between spin-split 2D hole states may also occur in this regime, as follows from the consideration[8] of hole spectra in tunnel-coupled quantum wells. These results are interesting in connection with the demonstration of the double electron layer tunneling transistor.[9]

The momentum-conserving tunneling of electrons in the systems with spin-orbit interaction has been proposed for usage in spin polarizers.[6,7,10] The spin polarization of electrons in such systems is determined by the direction of momentum **p**, and the electrons which tunnel with a given momentum are spin-polarized. However, due to the isotropy of the electron spectrum in the absence of the magnetic field, the total (averaged over the directions of **p**) tunneling current is not spin polarized. An application of the magnetic field parallel to the layers dramatically changes the situation, allowing one to select the electrons with given momenta. In this field, the electron dispersion laws in each quantum well become anisotropic due to mixing between spin-orbit and Pauli interactions.[11] More important, the diamagnetic shift of the 2D electron spectra[4,5] occurs, see Fig. 1(b). This leads to the appearance of *spin-polarized tunneling current*.

In this paper we consider the effect of in-plane magnetic field on the tunneling current between 2D electron layers in double quantum well systems (DQWs) with spin-orbit splitting of energy spectrum. The aim of the paper is to calculate both the absolute value of this current and its spin polarization as functions of the magnetic field, energy shift between the 2D levels, and applied voltage. We also investigate relative effects of the Zeeman splitting and diamagnetic shift on the spin-polarized tunneling current. By considering interaction of electrons with short-range and long-range static potentials, we study the influence of homogeneous and inhomogeneous broadening mechanisms on the tunneling current and its spin polarization.

The paper is organized as follows. In Sec. II we give the basic relation and derive the expression for the spin-polarized tunneling current. In Sec. III we present some analytical results obtained from this expression in the case of homogeneous broadening and show the results of numerical calculations. The conclusions are given in Sec. IV.

II. BASIC RELATIONS

Taking into account the interwell tunnel coupling described by the spin-independent tunneling matrix element, T, we use below the one-electron Hamiltonian,

$$\left| \begin{array}{cc} \hat{H}_u & T \\ T & \hat{H}_l \end{array} \right| . \tag{1}$$

The spin-dependent motion in jth QW is described by the Hamiltonians,

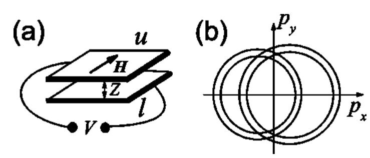

FIG. 1. (a) Schematic picture of the independently contacted DQW structure under an in-plane magnetic field $\mathbf{H} \| OY$. (b) Diamagnetic shift of the isoenergetic curves for DQWs with spin-split energy spectra.

$$\hat{H}_l = -\Delta/2 + \varepsilon_{\hat{\mathbf{p}}+\mathbf{p}_H/2} + \mathbf{w}_{l\hat{\mathbf{p}}} \cdot \hat{\boldsymbol{\sigma}} + V_{l\mathbf{x}},$$

$$\hat{H}_u = \Delta/2 + \varepsilon_{\hat{\mathbf{p}}-\mathbf{p}_H/2} + \mathbf{w}_{u\hat{\mathbf{p}}} \cdot \hat{\boldsymbol{\sigma}} + V_{u\mathbf{x}}, \tag{2}$$

where Δ is the energy shift between the 2D levels, which gives the level splitting without tunneling, $\varepsilon_{\mathbf{p}}=p^2/2m$ is the kinetic energy, $\mathbf{p}_H=(p_H,0,0)$, $p_H=eHZ/c$ is the characteristic momentum due to the magnetic field $\mathbf{H}\|OY$, $\hat{\boldsymbol{\sigma}}$ are the Pauli matrices, and $V_{u,l\mathbf{x}}$ are the potentials in u,l QWs. The vectors $\mathbf{w}_{u\mathbf{p}}=[(\mathbf{p}-\mathbf{p}_H/2)\times\mathbf{v}_u]+\mathbf{w}_H$ and $\mathbf{w}_{l\mathbf{p}}=[(\mathbf{p}+\mathbf{p}_H/2)\times\mathbf{v}_l]+\mathbf{w}_H$ contain the characteristic spin velocities in jth QW, $\mathbf{v}_j=(0,0,v_j)$, as well as the Zeeman energy, $\mathbf{w}_H=g\mu_B\mathbf{H}/2$, where g is the effective g factor and μ_B is the Bohr magneton.

Introducing the spin projection operator on the $\mathbf{e}$ direction, $\hat{P}_{\mathbf{e}}=[1+\mathbf{e}\cdot\hat{\boldsymbol{\sigma}}]/2$, we write the density matrix of electrons with spin along $\mathbf{e}$ as an anticommutator $\hat{\rho}_{\mathbf{e}}=(1/2)\times[\hat{P}_{\mathbf{e}},\hat{\rho}]_+$, where

$$\hat{\rho} = \begin{vmatrix} \hat{\rho}_u & \tilde{\rho} \\ \tilde{\rho}^+ & \hat{\rho}_l \end{vmatrix}$$

is written as a matrix in the basis of the layers u and l. The system of equations for the spinor density matrices in u and l layers, defined as $\hat{\rho}_{j\mathbf{e}}=(1/2)[\hat{P}_{\mathbf{e}},\hat{\rho}_j]_+$, is given by (see Refs. 5 and 12):

$$\frac{\partial \hat{\rho}_{u\mathbf{e}}}{\partial t} + \frac{i}{2\hbar}[\hat{P}_{\mathbf{e}},[\hat{H}_u,\hat{\rho}_u]]_+ = \frac{iT}{\hbar}(\tilde{\rho}_{\mathbf{e}} - \tilde{\rho}_{\mathbf{e}}^+),$$

$$\frac{\partial \hat{\rho}_{l\mathbf{e}}}{\partial t} + \frac{i}{2\hbar}[\hat{P}_{\mathbf{e}},[\hat{H}_l,\hat{\rho}_l]]_+ = -\frac{iT}{\hbar}(\tilde{\rho}_{\mathbf{e}} - \tilde{\rho}_{\mathbf{e}}^+), \tag{3}$$

where $\tilde{\rho}_{\mathbf{e}}=(1/2)[\hat{P}_{\mathbf{e}},\tilde{\rho}]_+$, and the nondiagonal part of the density matrix is written as

$$\tilde{\rho} = \frac{iT}{\hbar}\int_{-\infty}^{0} dt\, e^{\delta t} e^{i\hat{H}_u t/\hbar}(\hat{\rho}_u - \hat{\rho}_l)e^{-i\hat{H}_l t/\hbar}, \tag{4}$$

with $\delta \to +0$. Introducing the density of electrons in the jth layer with spin in the $\mathbf{e}$ direction as $n_{j\mathbf{e}}=\mathrm{Tr}\ \hat{\rho}_{j\mathbf{e}}$, we obtain the balance equations in the form

$$\frac{\partial n_{u\mathbf{e}}}{\partial t} - \frac{1}{\hbar}\mathrm{Tr}\hat{\rho}_u\hat{\boldsymbol{\sigma}}\cdot[\mathbf{e}\times\mathbf{w}_{u\hat{\mathbf{p}}}] = \frac{iT}{\hbar}\mathrm{Tr}(\tilde{\rho}_{\mathbf{e}} - \tilde{\rho}_{\mathbf{e}}^+),$$

$$\frac{\partial n_{l\mathbf{e}}}{\partial t} - \frac{1}{\hbar}\mathrm{Tr}\hat{\rho}_l\hat{\boldsymbol{\sigma}}\cdot[\mathbf{e}\times\mathbf{w}_{l\hat{\mathbf{p}}}] = -\frac{iT}{\hbar}\mathrm{Tr}(\tilde{\rho}_{\mathbf{e}} - \tilde{\rho}_{\mathbf{e}}^+). \tag{5}$$

The right-hand sides of these equations describe the interlayer tunneling, while the second terms on the left-hand sides, originating from the commutators $[\hat{H}_j,\hat{\rho}_j]$, describe free precession of the spins due to both spin-orbit and Pauli interactions. Since the longitudinal currents in strongly doped quantum wells are negligible, these last contributions exist only in the presence of a magnetic field. Once the spin quantization axis is chosen along $\mathbf{H}$, as $\mathbf{e}=\mathbf{h}\equiv\mathbf{H}/H$ (or $\mathbf{e}=-\mathbf{h}$), the spin precession in the layers is absent, and the second terms on the left-hand sides of Eq. (5) are zeros. This property can be demonstrated directly, by using the definitions of $\mathbf{w}_{j\hat{\mathbf{p}}}$ and taking into account the symmetry of the energy spectrum with respect to the y component of the momentum. The evolution of electron density in this case is determined only by electron transfer between the layers due to the tunneling, and the spin-polarized tunneling current can be written as $J_{\mathbf{h}}=e(\partial n_{u\mathbf{h}}/\partial t)=-e(\partial n_{l\mathbf{h}}/\partial t)$. Using Eq. (5) and the definition of $\tilde{\rho}_{\mathbf{e}}$, we present this current in the following form:

$$J_{\mathbf{h}} = \frac{ieT}{2\hbar L^2}\left\langle\left\langle \sum_{\lambda} (l\lambda|[\hat{P}_{\mathbf{h}},(\tilde{\rho}-\tilde{\rho}^+)]_+|l\lambda)\right\rangle\right\rangle, \tag{6}$$

where $\langle\langle\cdots\rangle\rangle$ is the averaging over random potentials and $|j\lambda)$ are the eigenstates of the Hamiltonian $\hat{H}_j$. Equation (6) is written in the T^2 approximation, when all contributions to $J_{\mathbf{h}}$ containing higher powers of T are neglected. Therefore, since $\tilde{\rho}$ is already proportional to T, see Eq. (4), we use the basis of single-layer states, $|l\lambda)$, to express the trace in Eq. (6) [one may equivalently use the basis $|u\lambda)$]. For the same reason, one should neglect the effect of tunnel coupling on $\hat{\rho}_u$ and $\hat{\rho}_l$ in the expression (4) when substituting the latter in Eq. (6). Doing it this way and introducing the eigenvalues $\varepsilon_{j\lambda}$ of the problem $\hat{H}_j|j\lambda)=\varepsilon_{j\lambda}|j\lambda)$, we transform Eq. (6) as follows:

$$J_{\mathbf{h}} = \frac{eT^2}{\hbar L^2}\left\langle\left\langle \sum_{\lambda\lambda'}\left[\frac{(u\lambda|\hat{P}_{\mathbf{h}}|l\lambda')(l\lambda'|u\lambda)}{i(\varepsilon_{l\lambda'}-\varepsilon_{u\lambda})+\delta} + \frac{(u\lambda|l\lambda')(l\lambda'|\hat{P}_{\mathbf{h}}|u\lambda)}{i(\varepsilon_{u\lambda}-\varepsilon_{l\lambda'})+\delta}\right](f_{l\lambda'}-f_{u\lambda})\right\rangle\right\rangle. \tag{7}$$

The distributions of electrons in (j,λ) states are supposed to be quasiequilibrium, $f_{j\lambda}=(j\lambda|\hat{\rho}_j|j\lambda)=f^{(j)}_{\varepsilon_{j\lambda}}$, with different chemical potentials in each QW determined by the doping level and transverse voltage, so that we further transform Eq. (7) into

$$J_{\mathbf{h}} = \frac{ieT^2}{\hbar L^2}\int d\varepsilon\langle\langle\mathrm{Tr}\hat{P}_{\mathbf{h}}\{f^{(l)}_{\varepsilon}(\hat{\mathcal{G}}^R_{u\varepsilon}\hat{\mathcal{A}}_{l\varepsilon} - \hat{\mathcal{A}}_{l\varepsilon}\hat{\mathcal{G}}^A_{u\varepsilon}) - f^{(u)}_{\varepsilon}(\hat{\mathcal{G}}^R_{l\varepsilon}\hat{\mathcal{A}}_{u\varepsilon} - \hat{\mathcal{A}}_{u\varepsilon}\hat{\mathcal{G}}^A_{l\varepsilon})\}\rangle\rangle, \tag{8}$$

where $\hat{\mathcal{G}}^R_{j\varepsilon}$ and $\hat{\mathcal{G}}^A_{j\varepsilon}$ are the retarded and advanced Green's function in the operator form, and $\hat{\mathcal{A}}_{j\varepsilon}=(\hat{\mathcal{G}}^A_{j\varepsilon}-\hat{\mathcal{G}}^R_{j\varepsilon})/2\pi i$ is the spectral density operator. The operator Green's functions, which satisfy the equation $(\varepsilon-\hat{H}_j\pm i\delta)\hat{\mathcal{G}}^{R,A}_{j\varepsilon}=\hat{1}$, can be viewed as matrices in both configuration and spin space and the trace Tr in Eq. (8) is taken over all coordinate and spin variables.

Below we assume that the potentials $V_{j\mathbf{x}}$ contain both short-range contribution and large scale, classically smooth contribution. Carrying out the averaging over the short-range contributions in Eq. (8), we imply that these contributions for different wells are statistically independent. Therefore, we obtain

$$J_{\mathbf{h}} = \frac{ieT^2}{\hbar L^2}\int d\varepsilon \sum_{\mathbf{p}} \mathrm{tr}_\sigma\{\hat{P}_{\mathbf{h}}\langle f_\varepsilon^{(l)}[\hat{G}_{u\varepsilon}^R(\mathbf{p},\mathbf{x})\hat{A}_{l\varepsilon}(\mathbf{p},\mathbf{x}) - \hat{A}_{l\varepsilon}(\mathbf{p},\mathbf{x})\hat{G}_{u\varepsilon}^A(\mathbf{p},\mathbf{x})] - f_\varepsilon^{(u)}[\hat{G}_{l\varepsilon}^R(\mathbf{p},\mathbf{x})\hat{A}_{u\varepsilon}(\mathbf{p},\mathbf{x}) - \hat{A}_{u\varepsilon}(\mathbf{p},\mathbf{x})\hat{G}_{l\varepsilon}^A(\mathbf{p},\mathbf{x})]\rangle\}, \quad (9)$$

where $\hat{G}_{j\varepsilon}^{R,A}(\mathbf{p},\mathbf{x})$ and $\hat{A}_{j\varepsilon}(\mathbf{p},\mathbf{x}) = \mathrm{Im}\,\hat{G}_{j\varepsilon}^R(\mathbf{p},\mathbf{x})/\pi$ are the Green's functions and spectral density functions in the Wigner representation. These functions are 2×2 matrices and the remaining trace tr_σ is taken over spin variable only. The averaging over the large-scale part of the random potential remains in Eq. (9), and it is denoted by $\langle\cdots\rangle$. The retarded Green's function satisfies the matrix Dyson equation

$$[\varepsilon - \hat{h}_{j\mathbf{p}} - U_{j\mathbf{x}} - \hat{\Sigma}_j^R]\hat{G}_{j\varepsilon}^R(\mathbf{p},\mathbf{x}) = \hat{1}, \quad (10)$$

where $\hat{h}_{u\mathbf{p}} = \Delta/2 + \varepsilon_{\mathbf{p}-\mathbf{p}_H/2} + \mathbf{w}_{u\mathbf{p}}\cdot\hat{\boldsymbol{\sigma}}$ and $\hat{h}_{l\mathbf{p}} = -\Delta/2 + \varepsilon_{\mathbf{p}+\mathbf{p}_H/2} + \mathbf{w}_{l\mathbf{p}}\cdot\hat{\boldsymbol{\sigma}}$ are the free-electron Hamiltonians in the momentum representation, $U_{j\mathbf{x}}$ is the large-scale part of random potential, and $\hat{1}$ is the unit matrix in the spinor basis. The self-energy function arising from the short-range scattering, $\hat{\Sigma}_j^R$, does not depend on ε, $\mathbf{p}$, and $\mathbf{x}$ for the case of scattering by zero-radius centers. Below we neglect the renormalization of $\hat{h}_{j\mathbf{p}}$ due to real part of $\hat{\Sigma}$ and express $\hat{\Sigma}_j^R$ through the energies of homogeneous broadening, $\hat{\Sigma}_j^R \simeq -i\gamma_j\hat{1}$. This approximation is valid in the case of weak spin splitting, under condition $|v_j|\overline{p} \ll \overline{\varepsilon}$, where $\overline{p}$ and $\overline{\varepsilon}$ are characteristic momentum and energy.

The expression (9) for the tunneling current can be considerably simplified if we take into account that, according to Eq. (10), $\hat{G}_R^{j\varepsilon}(\mathbf{p},\mathbf{x})$ does not contain the contributions proportional to the Pauli matrix $\hat{\sigma}_z$, because $\hat{h}_{j\mathbf{p}}$ does not contain such contributions. Therefore, the commutators of the Green's functions are proportional to $\hat{\sigma}_z$ and their trace with $\hat{P}_{\mathbf{h}}$ is zero. Using this property, one can carry out the permutations of the Green's functions under the trace in Eq. (9) so that the latter is rewritten as

$$J_{\mathbf{h}} = \frac{2\pi eT^2}{\hbar}\int d\varepsilon (f_\varepsilon^{(l)} - f_\varepsilon^{(u)})\int \frac{d\mathbf{p}}{(2\pi\hbar)^2} \times \langle \mathrm{tr}_\sigma\{\hat{P}_{\mathbf{h}}\hat{A}_{u\varepsilon}(\mathbf{p},\mathbf{x})\hat{A}_{l\varepsilon}(\mathbf{p},\mathbf{x})\}\rangle. \quad (11)$$

Finally, we assume that the large-scale potentials $U_{u\mathbf{x}}$ and $U_{l\mathbf{x}}$ are statistically independent. Thus, the tunneling current (11) can be written through the completely averaged spectral functions $\hat{A}_{j\varepsilon}(\mathbf{p}) = \langle\hat{A}_{j\varepsilon}(\mathbf{p},\mathbf{x})\rangle$ given by[12]

$$\hat{A}_{j\varepsilon}(\mathbf{p}) = \int_{-\infty}^0 \frac{dt}{2\pi\hbar} e^{\gamma_j t/\hbar - (\Gamma_j t/2\hbar)^2} e^{i(\hat{h}_{j\mathbf{p}}-\varepsilon)t/\hbar} + \mathrm{H.c.}, \quad (12)$$

where the inhomogeneous broadening in jth QW is determined in the quasiclassical approximation by the variance of the potential, $\Gamma_j = \sqrt{\langle U_{j\mathbf{x}}^2\rangle}$.

Substituting the expressions (10) into Eq. (11) and taking the spin trace $\mathrm{tr}_\sigma\ldots$, we finally obtain $J_{\mathbf{h}}$ in the form

$$J_{\mathbf{h}} = \frac{2\pi eT^2}{\hbar}\int d\varepsilon (f_\varepsilon^{(l)} - f_\varepsilon^{(u)})\int \frac{d\mathbf{p}}{(2\pi\hbar)^2}\Bigg\{A_{l\varepsilon}^{(+)}(\mathbf{p})A_{u\varepsilon}^{(+)}(\mathbf{p}) + \frac{(\mathbf{w}_{l\mathbf{p}}\cdot\mathbf{w}_{u\mathbf{p}})}{w_{l\mathbf{p}}w_{u\mathbf{p}}}A_{l\varepsilon}^{(-)}(\mathbf{p})A_{u\varepsilon}^{(-)}(\mathbf{p}) - \frac{(\mathbf{h}\cdot\mathbf{w}_{u\mathbf{p}})}{w_{u\mathbf{p}}}A_{l\varepsilon}^{(+)}(\mathbf{p})A_{u\varepsilon}^{(-)}(\mathbf{p}) - \frac{(\mathbf{h}\cdot\mathbf{w}_{l\mathbf{p}})}{w_{l\mathbf{p}}}A_{l\varepsilon}^{(-)}(\mathbf{p})A_{u\varepsilon}^{(+)}(\mathbf{p})\Bigg\}, \quad (13)$$

where $w_{j\mathbf{p}} = |\mathbf{w}_{j\mathbf{p}}|$ and the scalar functions $A_{j\varepsilon}^{(\pm)}$ are introduced according to

$$A_{j\varepsilon}^{(\pm)} = \int_{-\infty}^0 \frac{dt}{2\pi\hbar} e^{\gamma_j t/\hbar - (\Gamma_j t/2\hbar)^2} \times \left[\cos\frac{(\varepsilon_{j\mathbf{p}-} - \varepsilon)t}{\hbar} \pm \cos\frac{(\varepsilon_{j\mathbf{p}+} - \varepsilon)t}{\hbar}\right]. \quad (14)$$

In this equation we have used the dispersion laws for spin-split states, $\varepsilon_{j\mathbf{p}\pm} = \varepsilon_{j\mathbf{p}} \pm w_{j\mathbf{p}}$, where $\varepsilon_{l\mathbf{p}} = \varepsilon_{\mathbf{p}+\mathbf{p}_H/2} - \Delta/2$ and $\varepsilon_{u\mathbf{p}} = \varepsilon_{\mathbf{p}-\mathbf{p}_H/2} + \Delta/2$.

III. RESULTS

Below we consider the total tunneling current, $J_\perp = J_{\mathbf{h}} + J_{-\mathbf{h}}$, and the spin-polarized contribution, $\Delta J = J_{\mathbf{h}} - J_{-\mathbf{h}}$. Instead of ΔJ, it is convenient to introduce the degree of spin polarization, $S = \Delta J/J_\perp$. In the linear regime, the tunneling is characterized by the tunneling conductance $G = J_\perp/V$. We first assume that the temperature is low enough and replace the distribution functions $f_\varepsilon^{(j)}$ by the steplike functions $\theta(\varepsilon_{Fj} - \varepsilon)$, where quasi-Fermi energies are given by $\varepsilon_{Fl} = \varepsilon_F + eV/2$ and $\varepsilon_{Fu} = \varepsilon_F - eV/2$, where ε_F is the equilibrium Fermi energy. The effect of temperature appears to be not essential as soon as the temperature is small in comparison to the Fermi energy, see discussion in Sec. IV. Before presenting the results of numerical calculations of the total current and degree of spin polarization, we give some analytical results related to the case of homogeneous broadening, when Eqs. (13) and (14) lead to the tunneling conductance and spin polarization in the form

$$\begin{pmatrix} G \\ \Delta G \end{pmatrix} = \frac{e^2T^2}{4\pi\hbar^3}\int d\mathbf{p}\begin{pmatrix} \Phi(\mathbf{p}) \\ \Psi(\mathbf{p}) \end{pmatrix}, \quad S = \frac{\Delta G}{G}. \quad (15)$$

In these equations

$$\Phi(\mathbf{p}) = (\delta_{\gamma l+}\delta_{\gamma u+} + \delta_{\gamma l-}\delta_{\gamma u-})\left[1 + \frac{\mathbf{w}_{l\mathbf{p}}\cdot\mathbf{w}_{u\mathbf{p}}}{w_{l\mathbf{p}}w_{u\mathbf{p}}}\right] + (\delta_{\gamma l+}\delta_{\gamma u-} + \delta_{\gamma l-}\delta_{\gamma u+})\left[1 - \frac{\mathbf{w}_{l\mathbf{p}}\cdot\mathbf{w}_{u\mathbf{p}}}{w_{l\mathbf{p}}w_{u\mathbf{p}}}\right] \quad (16)$$

and

$$\Psi(\mathbf{p}) = (\delta_{\gamma l+}\delta_{\gamma u+} - \delta_{\gamma l-}\delta_{\gamma u-})\mathbf{h}\cdot\left[\frac{\mathbf{w}_{l\mathbf{p}}}{w_{l\mathbf{p}}} + \frac{\mathbf{w}_{u\mathbf{p}}}{w_{u\mathbf{p}}}\right] + (\delta_{\gamma l+}\delta_{\gamma u-} - \delta_{\gamma l-}\delta_{\gamma u+})\mathbf{h}\cdot\left[\frac{\mathbf{w}_{l\mathbf{p}}}{w_{l\mathbf{p}}} - \frac{\mathbf{w}_{u\mathbf{p}}}{w_{u\mathbf{p}}}\right], \quad (17)$$

where the shortcuts $\delta_{\gamma j\pm}$ stand for the Lorentz functions

$\pi^{-1}\gamma_j/[\gamma_j^2+(\varepsilon-\varepsilon_{j\mathbf{p}\pm})^2]$ and $\varepsilon=\varepsilon_F$. In the following, the broadening is assumed to be symmetric, $\gamma_l=\gamma_u=\gamma$.

The integrals over 2D momenta in Eq. (15) can be calculated analytically under the assumption $(j=l,u)$,

$$2m|v_j| \ll p_F, \tag{18}$$

where $p_F=\sqrt{2m\varepsilon_F}$ is the characteristic Fermi momentum. The condition (18) means that the spin velocities are small in comparison with the Fermi velocity. It is valid for any doped quantum wells with spin-orbit splitting of the energy spectrum. For the sake of simplicity, we neglect the Zeeman splitting term w_H. The calculation leads to the following results valid for arbitrary magnetic fields:

$$G=\frac{e^2T^2m}{2\pi\hbar^3}\sum_{\sigma_l,\sigma_u=\pm1}\left(1-\sigma_l\sigma_u\frac{p_F^2-p_H^2/2}{\sqrt{p_F^4-(m\Delta)^2}}\right)\times \mathrm{Im}\frac{1}{\sqrt{E_{\sigma_l\sigma_u}^2-p_H^2(p_F^2-p_H^2/4)/m^2}} \tag{19}$$

and

$$\Delta G=\frac{e^2T^2m^2}{2\pi\hbar^3 p_H}\sum_{\sigma_l,\sigma_u=\pm1}\times\left\{\frac{\sigma_l}{\sqrt{p_F^2+m\Delta}}\mathrm{Im}\frac{E_{\sigma_l\sigma_u}-p_H^2/2m}{\sqrt{E_{\sigma_l\sigma_u}^2-p_H^2(p_F^2-p_H^2/4)/m^2}}-\frac{\sigma_u}{\sqrt{p_F^2-m\Delta}}\mathrm{Im}\frac{E_{\sigma_l\sigma_u}+p_H^2/2m}{\sqrt{E_{\sigma_l\sigma_u}^2-p_H^2(p_F^2-p_H^2/4)/m^2}}\right\}. \tag{20}$$

The complex energy $E_{\sigma_l\sigma_u}$ is introduced according to

$$E_{\sigma_l\sigma_u}=(\sigma_l v_l+\sigma_u v_u)p_F+\Delta-2i\gamma. \tag{21}$$

The calculation of the square root from the complex expression in Eqs. (19) and (20) must be done under condition that $\mathrm{Im}\sqrt{\ldots}$ is negative. For example, in the absence of the magnetic field $(p_H=0)$ this rule means that the square root is equal to $E_{\sigma_l\sigma_u}$, and we recover the result of Ref. 7 for the tunneling conductance. The spin polarization in these conditions is zero.

To demonstrate the effect of magnetic field, we consider a symmetric structure, when $\Delta=0$ and spin velocities are equal in absolute value and have different signs, $v_l=-v_u=v$. The different signs of the spin velocities in the two layers is physically understandable because the directions of the potential gradients in the l and u wells of the double-well structure are opposite to each other. The case $\Delta=0$ at $|v_l|=|v_u|$ corresponds to matched 2D electron densities in the wells, $n_l=n_u=n=(p_F/\hbar)^2/2\pi$. In these conditions, the tunneling current in the absence of the scattering is zero at $H=0$, because the states with the same direction of spin are out of resonance. With the increase of the magnetic field, when p_H reaches $2m|v|$, see Fig. 2, the tunneling conductance has a peak corresponding to the resonance of the states with the same direction of spin. The spin polarization in these conditions has a maximum. The tunneling at $p_H>2m|v|$ is described by a simple expression derived from Eqs. (19)–(21) under the assumptions $\gamma\to0$ and $p_H\ll p_F$,

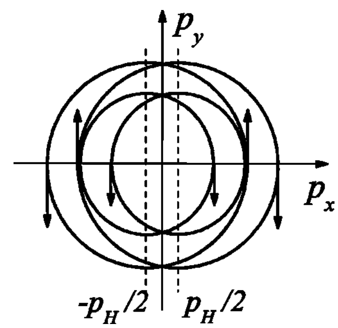

FIG. 2. Fermi surfaces for symmetric DQWs under condition of the maximum spin polarization of the tunneling current. The arrows show spin orientation at $p_y=0$ for each branch of the energy spectrum.

$$G=\frac{2e^2T^2m}{\pi\hbar^3 p_F\sqrt{(p_H/m)^2-(2v)^2}},\quad S=\frac{2mv}{p_H}. \tag{22}$$

The conductance diverges at $p_H\simeq 2m|v|$ as $(p_H-2m|v|)^{-1/2}$. The spin polarization reaches 100% in the peak and slowly decreases with the increase of H. As p_H becomes smaller than $2m|v|$, the tunneling conductance rapidly decreases with the decrease of the magnetic field. To evaluate the conductance and the spin polarization in this weak-field region one should assume a finite broadening, $\gamma\neq0$.

Further, we consider the case when one quantum well, say u, is symmetric so that $v_u=0$, while the other one is not symmetric, $v_l=v\neq0$. The spin-orbit splitting of electron states in the u layer is absent and, for this reason, the spin relaxation there is considerably suppressed. In this case, the spin-polarized electrons injected to the u layer from the l layer can keep their polarization for a long time, determined only by spin-dependent scattering processes, which is important[13] for spintronics applications. We again neglect Zeeman splitting and assume $\Delta=0$, which, under the assumed condition (18), corresponds to matched electron densities in the wells. Under these conditions, and under assumptions $\gamma\to0$ and $p_H\ll p_F$, the conductance and spin polarization at $p_H>m|v|$ are given by the following expressions:

$$G=\frac{2e^2T^2m}{\pi\hbar^3 p_F\sqrt{(p_H/m)^2-v^2}},\quad S=\frac{mv}{p_H}, \tag{23}$$

which are very similar to those given by Eq. (22). The resonance peak of the conductance appears at $p_H=m|v|$, the spin polarization reaches there 100% and decreases with the increase of H.

It is important to mention that the neglect of Zeeman splitting in comparison to the spin-orbit splitting is well justified in the InAs(InGaAs) quantum wells at small enough magnetic fields. Indeed, there is a strong inequality,

$$w_H \ll p_H p_F/m, \tag{24}$$

which is field independent, since both w_H and p_H are proportional to H. Estimating the effective mass m as 0.04 of the

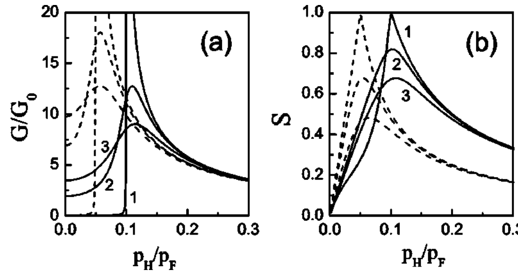

FIG. 3. Tunneling conductance (a) and spin polarization of the tunneling current (b) as functions of the magnetic field at $\Delta=0$ and $\gamma/\varepsilon_F=10^{-4}$ (1), 0.02 (2), and 0.04 (3). Solid lines, $r_l=-r_u=0.1$; dashed lines, $r_l=0.1$ and $r_u=0$.

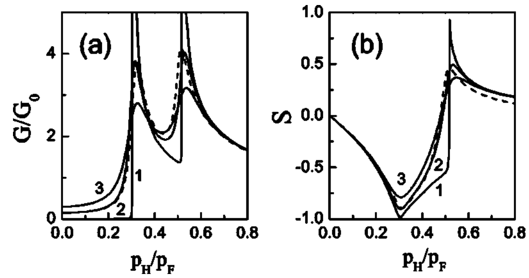

FIG. 4. Tunneling conductance (a) and spin polarization of the tunneling current (b) as functions of the magnetic field at $|\Delta|=0.8\varepsilon_F$ and $\gamma/\varepsilon_F=10^{-4}$ (1), 0.02 (2), and 0.04 (3). The dashed lines show the result of numerical calculation at $\gamma/\varepsilon_F=0.02$ taking into account Zeeman splitting.

free electron mass m_0, and considering typical electron densities, $n\simeq 10^{12}$ cm^{-2}, and typical interwell separation in the DQWs with independent contacts, $Z=25$ nm, one can find that the condition (24) is well satisfied despite of the large value of the g factor in InAs quantum wells,[14] $g\simeq -13$. On the other hand, as shown above, the resonance peak of the conductance and spin polarization at $\Delta=0$ exists around $p_H \simeq m|v_{l,u}|$. According to Eq. (24), in this region of magnetic fields the spin-orbit spitting of the energy spectrum near the Fermi level, estimated as $|v_{l,u}|p_F$, appears to be much greater than the Zeeman splitting w_H, and the latter should be neglected. At higher magnetic fields, a noticeable contribution of Zeeman splitting appears, but it does not modify the tunneling current and spin polarization considerably, as demonstrated below by numerical calculations.

The inverse-square-root divergencies of the tunneling conductance in Eqs. (22) and (23) are smeared in the presence of finite broadening. The broadening also decreases the degree of spin polarization. To investigate these effects, and to estimate the degree of spin polarization expected in realistic conditions, we plot below the dependence of the tunneling conductance and spin polarization given by Eqs. (19)–(21) on the magnetic field. The field is described by a dimensionless parameter p_H/p_F. The other relevant dimensionless parameters are the ratios of the spin velocities to the Fermi velocity, $r_j=2mv_j/p_F$, and the ratios of the level separation and broadening energies to the Fermi energy, Δ/ε_F and γ/ε_F, respectively. In Fig. 3 we plot G, in units of $G_0 = e^2T^2m/\pi\hbar^3\varepsilon_F$, and S at $\Delta=0$ for several values of γ/ε_F using $r_l=-r_u=0.1$ and $r_l=0.1$, $r_u=0$. At $n=10^{12}$ cm^{-2} and $m=0.04m_0$ the value $|r_{l,u}|=0.1$ corresponds to $|v_{l,u}|=3.6\times 10^6$ cm/s, which is larger than the typical value 1.5×10^6 cm/s obtained in experiments[15] on single InGaAs quantum wells, but smaller than the maximal value 4.85×10^6 cm/s reported recently.[16] Therefore, we consider r_l and r_u used in the calculations as reasonable. The given values of broadening energies are also reasonable. For example, $\gamma/\varepsilon_F=0.02$ for $n=10^{12}$ cm^{-2} and $m=0.04m_0$, when $\varepsilon_F \simeq 60$ meV, corresponds to $\gamma=1.2$ meV (or $\hbar/\tau=2\gamma=2.4$ meV, where τ is the scattering time of electrons).

The resonance tunneling peak in Fig. 3(a) corresponds to the maximum of the spin polarization in Fig. 3(b). According to the simple consideration given above, the peak occurs at $p_H/p_F\simeq r_l$ for $r_l=-r_u=0.1$ and at $p_H/p_F\simeq r_l/2$ for $r_l=0.1$ and $r_u=0$. In the regions of higher field, the behavior of G and S is close to that described by Eqs. (22) and (23). In the case of negligible broadening, the spin polarization reaches unity in a kinklike fashion. The broadening of the resonance tunneling peak is accompanied with the decrease of the maximum spin polarization, which, nevertheless, remains rather high (60–80%) for the physically reasonable values of broadening considered here. For obvious reasons, the inversion of the magnetic field (negative p_H) would lead to a symmetric transformation of G and to an antisymmetric transformation of S.

The field dependence of G and S at large level separation $|\Delta|=0.8\varepsilon_F$ is illustrated in Fig. 4, where we assumed $r_l=-r_u=0.1$ and used the same values of broadening energy γ as in Fig. 3. The characteristic feature of this dependence is the appearance of two peaks, corresponding to $p_H/p_F \simeq |\Delta|/2\varepsilon_F \pm r_l$. In the region between the peaks the sign of the spin polarization is changed (from -1 to 1 in the case of negligible broadening). The first peak of S occurs in a kinklike fashion, similar as in Fig. 3, while the second one is more narrow and more sensitive to the broadening. In Fig. 4, the results derived from Eqs. (19)–(21) are compared with the results of exact numerical calculation using Eqs. (15)–(17). These calculations take into account Zeeman splitting estimated for the parameters $n=10^{12}$ cm^{-2}, $Z=25$ nm, $m=0.04m_0$, and $g=-13$. One may notice that the difference becomes noticeable at $p_H/p_F>0.3$ (it is not visible for the region of fields used in Fig. 3). Nevertheless, this difference remains small and the qualitative behavior given by Eqs. (19)–(21) is not modified by the Pauli interaction effect.

In Fig. 5 we plot the results of numerical calculation of the tunneling conductance for the case of symmetric inhomogeneous broadening $\Gamma_l=\Gamma_u=\Gamma$ as well as for the realistic case when both broadening mechanisms are presented. The difference in the resonance tunneling peak shape for two broadening mechanisms reflects the different behavior of Lorentz and Gauss functions. A large degree of spin polarization is possible even for strong inhomogeneous broadening $\Gamma=0.1\varepsilon_F$. The homogeneous broadening suppresses the spin polarization more considerably than the inhomogeneous one.

So far we considered the linear regime of tunneling. Let us discuss the case of a finite voltage V applied between the 2D layers. To calculate the tunneling current, $J_\perp$, and its spin-polarized part, ΔJ, one may use expressions for G and

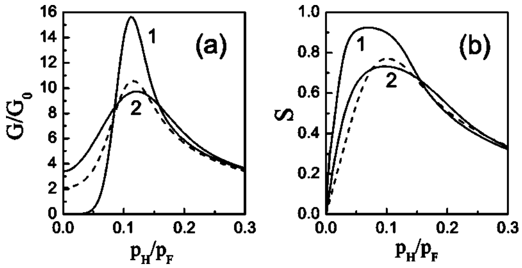

FIG. 5. Tunneling conductance (a) and spin polarization of the tunneling current (b) as functions of the magnetic field at $\Delta=0$. Solid lines, inhomogeneous broadening with $\Gamma=0.04\ \varepsilon_F$ (1) and $\Gamma=0.1\ \varepsilon_F$ (2). Dashed lines show the case of mixed broadening with $\gamma=0.02\ \varepsilon_F$ and $\Gamma=0.04\ \varepsilon_F$.

ΔG as functions of energy ε [for example, formally replacing p_F in Eqs. (19)–(23) by $\sqrt{2m\varepsilon}$] and integrate these expressions over ε in the interval $[\varepsilon_F - eV/2, \varepsilon_F + eV/2]$. It is important that the application of a finite voltage changes the energy separation Δ. It also causes a redistribution of electron density over the layers and a renormalization of the spin velocities. These effects are considered in Ref. 7 for the simplest case of initially symmetric DQWs, with $\Delta=0$ and $v_u = -v_l$ at $V=0$. Below we use the results of Ref. 7 written as

$$\Delta = -\frac{eV}{1 + a_B/2Z},$$

$$v_l = v\left(1 - \frac{a_B}{a_B + 2Z}\frac{eV}{2\varepsilon_F}\right),$$

$$v_u = -v\left(1 + \frac{a_B}{a_B + 2Z}\frac{eV}{2\varepsilon_F}\right), \tag{25}$$

where $a_B = \hbar^2\epsilon/me^2$ is the Bohr radius, which characterizes the screening effects in 2D layers.

The magnetic-field dependence of the nonlinear tunneling conductance $J_\perp/V$ calculated in the way described above is plotted in Fig. 6 for several values of the dimensionless bias eV/ε_F. When using Eq. (25), we set $2mv/p_F = 0.1$, $Z = 25$ nm, and estimated the Bohr radius for $m = 0.04m_0$ and $\varepsilon = 12$. The modifications of the tunneling current and its spin polarization are caused mostly by the dependence of Δ on V. The effects of finite voltage drop and renormalization of the spin velocities are far less essential. The first effect is not important because the energy dependence of G and ΔG does not have any peculiarities near the Fermi energy [this is seen in the most obvious way from the simple expressions (22) and (23)], while the second one can be neglected because the main contribution to G and ΔG of Eqs. (19) and (20) comes from the terms containing the quantity $v_l - v_u$, which is not renormalized by the applied bias, see Eq. (25).

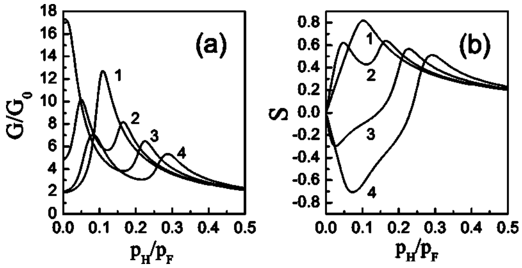

FIG. 6. Nonlinear tunneling conductance (a) and spin polarization of the tunneling current (b) as functions of the magnetic field for the case of homogeneous broadening with $\gamma=0.02\ \varepsilon_F$ at $eV/\varepsilon_F = 0, 0.15, 0.3$, and 0.45 (curves 1–4, respectively).

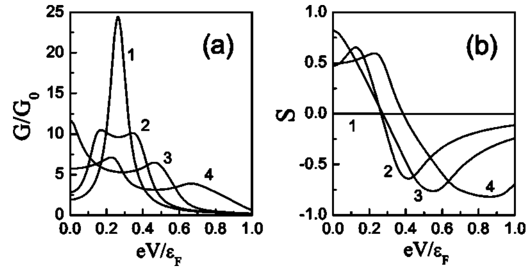

FIG. 7. Nonlinear tunneling conductance (a) and spin polarization of the tunneling current (b) as functions of the applied voltage for the case of homogeneous broadening with $\gamma=0.02\ \varepsilon_F$ at $p_H/p_F = 0, 0.05, 0.1$, and 0.2 (curves 1–4, respectively).

The dependence of the nonlinear tunneling conductance on the applied bias for several values of the dimensionless magnetic field p_H/p_F is shown in Fig. 7. Again, due to the dependence of Δ on V, the nonohmic behavior is pronounced. The spin polarization changes its sign with the increase of V, but its absolute value can be fairly large even at high applied bias $eV \sim \varepsilon_F$.

IV. CONCLUDING REMARKS

We have investigated the effect of an in-plane magnetic field on the tunneling current between independently contacted quantum wells with spin splitting of the energy spectrum. The main result of our studies is that a considerable spin polarization of electrons in a 2D layer can be achieved by means of tunneling injection of electrons into this layer from the adjacent 2D layer in the presence of a weak magnetic field applied along the layers. The scale of this field, corresponding to the maximum spin polarization, is estimated from the relation $p_H = (eH/c)Z \sim m|v|$, where Z is the interlayer distance and v is the spin velocity determining the spin splitting of the energy spectrum. Using the typical values $|v| < 4\times10^6$ cm/s and $Z > 200$ nm, one can find that the required magnetic fields for InAs (InGaAs)-based structures are smaller than 0.5 T. For the DQWs of small size, such fields, in principle, can be created by ferromagnetic films deposited on the surface of the sample.[17]

By investigating the effects of homogeneous and inhomogeneous broadening of energy spectrum, we have found that although both these effects suppress the spin polarization, a fairly large degree of polarization (60–80%) can be achieved for realistic values of the broadening energies, including the case when both broadening mechanisms are present. We have shown that the inhomogeneous broadening does not suppress the spin polarization considerably. In addition, the resonance peak of spin polarization for this mechanism appears to be wider than in the case of homogeneous broadening, which may offer certain advantages for applications. The conclusion about the weaker role of inhomogeneous broadening is encouraging since it is this broadening that is important in InGaAs structures because of considerable large-scale interface roughness.

Let us discuss the effect of electron temperature T_e on the tunneling conductance. To take it into account, one should consider G and ΔG as functions of energy, as described in Sec. III, and integrate them over the energy with the weight factor $-(\partial f_\varepsilon / \partial \varepsilon)$. Because of smooth energy dependence of G and ΔG discussed above, the relative corrections due to finite temperature appear to be of the order $(T_e/\varepsilon_F)^2$. Thus, both the tunneling conductance and spin polarization are not sensitive to the temperature unless the latter is comparable to the Fermi energy. The condition $(T_e/\varepsilon_F)^2 \ll 1$ is easy to achieve even at $T_e = 77$ K in the InAs-based structures, where the 2D electron densities are high and the Fermi energy typically exceeds 50 meV. However, with the increase of temperature, an additional broadening mechanism due to scattering of electrons by phonons becomes important, and this can lead to additional suppression of the spin polarization.

The main difficulty for application of the tunneling injection of spin-polarized 2D electrons is the small value of the tunneling current. One of the ways to overcome it is to apply high voltages between the layers. Our calculations show that, by a proper choice of the value of in-plane magnetic field, the degree of spin polarization can be made high (60–80%) even when the applied bias eV is comparable to the Fermi energies in the layers. When it is so, the flux density of incoming spin-polarized electrons is roughly estimated as $\Delta J/e \sim G\varepsilon_F/e \sim 10T^2 m/\pi\hbar^3$ (the approximate numerical factor of 10 is taken from our calculations, see Figs. 3–7). By equating this quantity to the rate of spin relaxation per unit square of the 2D layer, n_s/τ_s, one can estimate the quasistationary density of spin-polarized electrons, n_s, provided that the spin relaxation time, τ_s, is known. Therefore, if the quantum wells are separated by thin enough tunneling barrier (so that T^2 is large), a considerable fraction of spin-polarized electrons in the 2D layers can appear as a result of the highly spin-polarized tunneling current studied in this paper. For example, using the spin relaxation time $\tau_s \simeq 10$ ps in InAs 2D layers,[18] the Fermi energy $\varepsilon_F \simeq 30$ meV (corresponding to a typical electron density $n = 5 \times 10^{11}$ cm^{-2}), and the tunneling matrix element $T = 0.2$ meV (attainable in InAs/GaAlSb/InAs double quantum well structures with tunnel barriers of about 7 nm thick), we obtain the relative spin polarization $n_s/n \simeq 0.2$. This value increases if T and τ_s are bigger and n is smaller. A consideration of the effect of such spin polarization in the framework of the T^2 approximation for the tunneling current and quasiequilibrium distribution functions in the layers is possible under conditions that τ_s is much greater than the momentum and energy relaxation times in the layers. For possible applications of independently contacted double quantum wells in spintronics, as spin filters, it is desirable to avoid the spin relaxation by creating the structures where the tunnel coupling between the layers occurs in a narrow lateral strip whose width is smaller than the spin diffusion length $\sqrt{D\tau_s}$ (D is the diffusion coefficient), but large in comparison to the ordinary diffusion length $\sqrt{D\tau}$ so that the formalism applied in our paper remains valid. In this case, the spin-polarized electrons come out from the tunnel-coupled region before the spin relaxation occurs, and one may use the strong spin polarization of the tunneling current directly.

In conclusion, our theoretical analysis demonstrates that the efficient injection of spin-polarized electrons can be achieved in double quantum well InAs(InGaAs) structures, where splitting of electron spectrum caused by the spin-orbit interaction is considerable. We hope that this result will stimulate technological efforts towards creation of independently contacted quantum wells based on these materials.

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