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Field-assisted electron transport through a symmetric double-well structure with spin–orbit coupling and the Fano-resonance induced spin filtering*

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We have investigated theoretically the field-driven electron-transport through a double-quantum-well semiconductor-heterostructure with spin–orbit coupling. The numerical results demonstrate that the transmission spectra are divided into two sets due to the bound-state level-splitting and each set contains two asymmetric resonance peaks which may be selectively suppressed by changing the difference in phase between two driving fields. When the phase difference changes from 0 to π, the dip of asymmetric resonance shifts from one side of resonance peak to the other side and the asymmetric Fano resonance degenerates into the symmetric Breit–Wigner resonance at a critical value of phase difference. Within a given range of incident electron energy, the spin polarization of transmission current is completely governed by the phase difference which may be used to realize the tunable spin filtering.

Keywords: Fano effect, Floquet channel, spin–orbit coupling
PACC: 7210, 7225, 2570E

1. Introduction

The ability to generate, manipulate and detect polarized spins in nanostructures not only can provide a powerful tool for investigating basic properties of spins in coherent electronic systems, but also is an indispensable basis guiding the design of spintronic devices in which the spin filter is an important building block. Thus much effort has been made to find the effective mechanism for achieving tunable spin filtering. The mechanisms commonly used for creating spin-polarized currents rely on a large Zeeman splitting in a strong magnetic field[1,2] or hybrid micromagnet-semiconductor structure.[3−5] However, the Zeeman splitting of spin states is not the most practical way to realize the spin filtering because the needed magnetic-field strengths are often too strong and on-chip placement of micromagnets is also required. Moreover, the spin injection efficiency from a ferromagnetic metallic contact into a semiconductor is small due to the conductivity mismatch.[6] The achievement of spin filtering by means of intrinsic spin-dependent effects in semiconductors is, therefore, highly desirable and also very intriguing from the point of view of fundamental science. Various designs based on semiconductor heterostructures can be used to create spin-polarized currents without ferromagnets or magnetic fields.[7−11] Recently Perel *et al*[12] and Glazov *et al*[13] have demonstrated that in the heterostructures made of semiconductor material with bulk inversion asymmetry, the k^3 Dresselhaus spin–orbit interaction, which couples spin states with space motion of the conduction electrons, can lead to a clear spin-dependent splitting of the transmission resonances through a double-barrier system. The applied oscillating field in a spin–orbit coupling region provides a feasible scheme of controlling the spin-dependent transport through the system.[14] In the present work, we investigate the photon-mediated resonance transmission through a symmetric semiconductor double-well structure with the Dresselhaus spin–orbit coupling. Two oscillating fields are applied correspondingly in the two potential-well regions. We propose a scheme to generate the spin filter based on

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the spin-resolved Fano resonance. Since the resonance tunnelling through a thin barrier between two wells leads to the bound-state level-splitting, corresponding transmission resonances are also divided into two sets and each set contains two asymmetric resonance peaks induced by the Dresselhaus spin–orbit coupling. The difference in phase between the two oscillating fields plays an important role in determining the resonance line phase and can be used to control the spin polarization of the transmission electrons.

The Fano-type resonance, which is induced by the interference effect between a localized state and the continuum band[15] of the system, has been known to be related to many different branches of physics and observed in a great variety of experiments including atomic photoionization,[16] electron and neutron scattering,[17] Raman scattering,[18] photoabsorption in quantum well structures[19,20] and electron transport through mesoscopic systems.[21−23] The Fano resonances of conductance have a line shape in the following form:[24−26]

$$G(\epsilon) = G_{\rm d}\frac{(\epsilon+q)^2}{\epsilon^2+1} \quad \text{with} \quad \epsilon = \frac{(E-E_{\rm R})}{\Gamma/2}, \qquad (1)$$

where $G_{\rm d}$ is the nonresonant conductance, $E_{\rm R}$ and Γ are the position and the width of Fano resonance, q is the asymmetry parameter (Fano factor). The Fano resonance is proportional to the ratio between transmission amplitudes of the resonant and the nonresonant channels and is a measure of coupling strength between the continuum state and the bound state.[25] A dip lies on the left-hand side of asymmetric resonant peak for $q > 0$ and on the right-hand side for $q < 0$, while the conductance displays a symmetric Breit–Wigner type resonance in the limit of $q \to \infty$. In the limit of $q \to 0$, the nonresonant transmission dominates, resulting in a symmetric dip. In our model, the asymmetric parameter q may be adjusted by the difference in phase between the two oscillating fields, which have an effect on the electron-transmission through the system and thus modify the line shape of the transmission resonance.

In Section 2, we introduce our model system and solve the time-periodic Schrödinger equation by using the Floquet theorem, and we obtain the spin-dependent transmission probability and conductance. In Section 3, we give the numerical results within the framework given in Section 2. Finally we make some concluding remarks in Section 4.

2. The model and formalism

We consider the transmission of a single-electron with an incident wave vector $k = \left(k_{||}, k_z\right)$ above a double potential wells composed of five layers of semiconductor InP/GaSb/InP/GaSb/InP. Two applied oscillating fields are confined in the two wells formed by a narrow band gap layer (GaSb) with a larger Dresselhaus spin–orbit coupling. The electron passes through five regions I, L, II, and R, III respectively from left to right (see Fig.1). The effective mass of electron out of well is μ_1. The capital letters L and R denote the

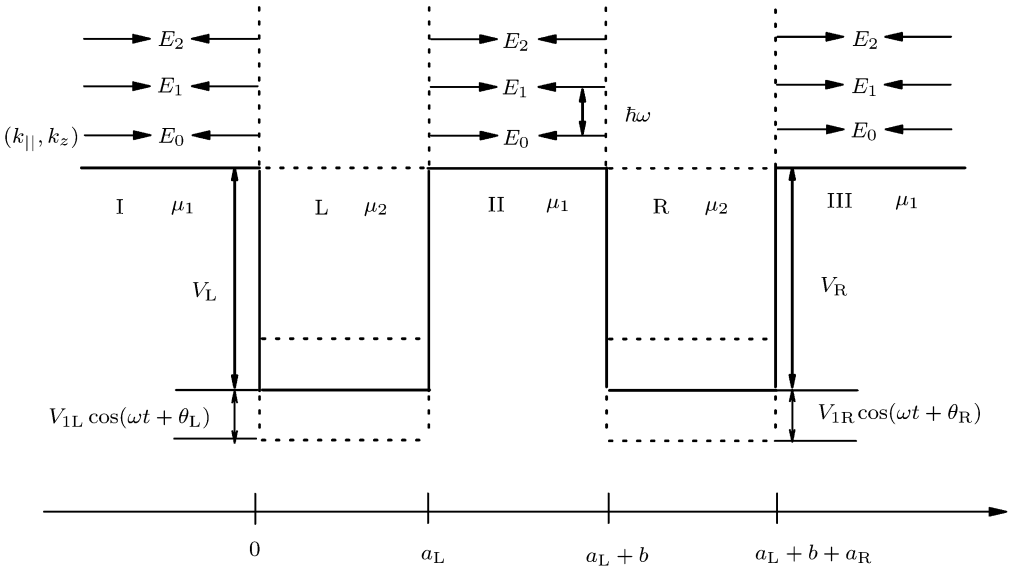


Fig.1. Model sketch of electron transport through a double-well structure with spin–orbit coupling and two applied oscillating fields.

left and the right wells with an electron effective mass μ_2 respectively. The time-periodic potential reads

$$V(z,t)=\begin{cases}-V_{\rm L}+V_{1\rm L}\cos(\omega t+\theta_{\rm L}), & 0\le z\le a_{\rm L},\\ -V_{\rm R}+V_{1\rm R}\cos(\omega t+\theta_{\rm R}), & a_{\rm L}+b\le z\le a_{\rm L}+b+a_{\rm R},\\ 0, & \text{otherwise},\end{cases}\tag{2}$$

where $V_{\rm L}$ and $V_{\rm R}$ are the static well depths and $V_{1\rm L}$, $V_{1\rm R}$ and ω denote the amplitudes and frequency of applied oscillating fields, respectively. k_z is the wave vector normal to the well plane and points to the direction of transmission. We assume that the temperature is so low that the electron–phonon interactions can be neglected. Thus the electron motion in each layer of the semiconductor structure may be described by a time-dependent Schrödinger equation

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\Phi(z,s,t)=\hat{H}\,\Phi(z,s,t)\tag{3}$$

with Hamiltonian being

$$\hat{H}=-\frac{\hbar^2}{2\mu}\frac{\partial^2}{\partial z^2}+\frac{\hbar^2k_{||}^2}{2\mu}+V(z,t)+\hat{H}_{\rm D},\tag{4}$$

where μ is the effective mass of electron and $\hat{H}_{\rm D}$ denotes the spin-dependent Dresselhaus term in Zinc-blende structure semiconductors.[27] In a quantum well of a finite depth, with the electron wavefunction confined in the growth direction, one should consider the wave vector k_z as an operator $-\mathrm{i}\partial/\partial z$. Assuming that the kinetic energy of incident electron is substantially smaller than the well depth $V_{\rm L,R}$, then the Dresselhaus term $\hat{H}_{\rm D}$ is simplified into[12,13]

$$\hat{H}_{\rm D}=\begin{cases}\gamma\left(\hat{\sigma}_xk_x-\hat{\sigma}_yk_y\right)\dfrac{\partial^2}{\partial z^2}, & 0\le z\le a_{\rm L},a_{\rm L}+b\le z\le a_{\rm L}+b+a_{\rm R},\\ 0 & \text{otherwise},\end{cases}\tag{5}$$

where γ is a material constant describing the strength of spin–orbit coupling, $\hat{\sigma}_i$ $(i=x,y)$ are the Pauli matrices. We have assumed that $\gamma\approx 0$ outside the wells. Inserting $\Phi(z,s,t)=\psi(z,t)\chi_{\rm s}$ with $\chi_{\rm s}$ being the spin wave function into Schrödinger Eq.(3) we obtain

$$\chi_{\pm}=\frac{1}{\sqrt{2}}\begin{pmatrix}1\\ \mp\mathrm{e}^{-\mathrm{i}\varphi}\end{pmatrix},\tag{6}$$

where $\chi_{\pm}$ are two pure spinor-states and φ denotes the polar angle of the wave vector $\boldsymbol{k}_{\|}=\left(k_{\|}\cos\varphi,k_{\|}\sin\varphi\right)$ in the XY plane. The orientation of electron spin $\boldsymbol{s}_{\pm}(\boldsymbol{k}_{||})=\chi_{\pm}^{+}\hat{\sigma}\chi_{\pm}=(\mp\cos\varphi,\pm\sin\varphi,0)$ corresponding to the eigenstates $\chi_{\pm}$ depends on the direction of wave vector $\boldsymbol{k}_{\|}$. In equilibrium the momentum distribution of incident electrons is isotropic in the XY plane and therefore the average spin of transmitted electrons vanishes.[12] In order to obtain the net spin polarization in transmitted electrons, we consider the case in which the in-plane wave vector of incident electron is along a fixed direction, for example $\varphi=0$, which may be achieved by applying an electric field along the $-X$ direction. In the spin subspace the Hamiltonian $\hat{H}$ reduces to

$$\hat{H}_{\pm}=\chi_{\pm}^{+}\hat{H}\chi_{\pm}=-\frac{\hbar^2}{2\mu_{\pm}}\frac{\partial^2}{\partial z^2}+\frac{\hbar^2k_{\|}^2}{2\mu_2}+V(z,t),\tag{7}$$

where the modified effective mass of electron depends not only on the Dresselhaus coupling constant γ and the in-plane electron wave vector $k_{||}$, but also on the orientation of electron spin, expressed as $\mu_{\pm}=\mu_2\left(1\pm\gamma 2\mu_2k_{\|}/\hbar^2\right)^{-1}$. Although only a little modification of electron effective mass is made, it plays an important role in generating the asymmetric Fano resonance peak splitting. Thus the Schrödinger equation can be rewritten as

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\Phi_{\pm}(z,t)=\hat{H}_{\pm}\,\Phi_{\pm}(z,t)\tag{8}$$

with

$$\Phi_{\pm}(z,t)=\chi_{\pm}\psi_{\pm}(z,t)\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}),\tag{9}$$

where $\boldsymbol{\rho}=(x,y)$ is a vector in the well-plane. Inserting Eqs.(7) and (9) into Eq.(8) we obtain the equation of the wavefunctions $\psi_{\pm}(z,t)$ as follows:

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\psi_{\pm}(z,t)=-\frac{\hbar^2}{2\mu_{\pm}}\frac{\partial^2}{\partial z^2}\psi_{\pm}(z,t)+\left(\frac{\hbar^2k_{\|}^2}{2\mu_2}+V(z,t)\right)\psi_{\pm}(z,t).\tag{10}$$

The potential $V(z,t)$ is a periodic function of time inside the potential well, thus according to the Floquet theorem,[28–30] Eq.(10) has the solution in the form $\psi_{\pm}(z,t)=\varphi_{\pm}(z,t)\exp(-\mathrm{i}E_{\rm F}^{\pm}t/\hbar)$, where $E_{\rm F}^{\pm}$ denote

the Floquet energy eigenvalues ranging from zero to $\hbar\omega$ and $\varphi_{\pm}(z,t)$ each are a periodic function of time: $\varphi_{\pm}(z,t)=\varphi_{\pm}(z,t+T)$ with a period (T) of $\frac{2\pi}{\omega}$. Substituting $\psi_{\pm}(z,t)$ into Eq.(10), the wavefunctions in the potential well are written as

$$\begin{aligned}\Phi_{\pm}^{\mathrm{L,R}}(z,t)=&\chi_{\pm}\sum_{n=-\infty}^{+\infty}\sum_{m=-\infty}^{+\infty}[a_m^{\pm}\mathrm{e}^{\mathrm{i}q_m^{\pm}z}+b_m^{\pm}\mathrm{e}^{-\mathrm{i}q_m^{\pm}z}]\\&\times J_{n-m}(\frac{V_{1\mathrm{L},1\mathrm{R}}}{\hbar\omega})\mathrm{e}^{-\mathrm{i}E_{zn}^{\pm}t/\hbar}\mathrm{e}^{-\mathrm{i}(n-m)\theta_{\mathrm{L,R}}}\\&\times\mathrm{e}^{-\mathrm{i}\frac{V_{1\mathrm{L},1\mathrm{R}}}{\hbar\omega}\sin(\theta_{\mathrm{L,R}})}\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}-\mathrm{i}E_{||}^{\pm}t/\hbar),\end{aligned}\tag{11}$$

where $J_n(x)$ is the Bessel function of the first kind, $a_m^{\pm}$ and $b_m^{\pm}$ are constant coefficients and

$$q_m^{\pm}=\left[\frac{2\mu_{\pm}}{\hbar^2}(E_{\mathrm{F}}^{\pm}+m\hbar\omega-\frac{\hbar^2k_{||}^2}{2\mu_2}+V_{\mathrm{L,R}})\right]^{1/2}$$

with m being an integer.

Since electrons incident to the oscillating region will be scattered inelastically into an infinite number of Floquet sidebands, the wavefunction outside the well can be written as the superposition of waves with all values of energy, i.e.

$$\begin{aligned}\Phi_{\pm}^{(i)}(z,t)=&\chi_{\pm}\sum_{n=-\infty}^{+\infty}[c_n^{(i)\pm}\mathrm{e}^{\mathrm{i}k_{zn}^{\pm}z}+d_n^{(i)\pm}\mathrm{e}^{-\mathrm{i}k_{zn}^{\pm}z}]\\&\times\mathrm{e}^{-\mathrm{i}E_{zn}^{\pm}t/\hbar}\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}-\mathrm{i}E_{||}^{\pm}t/\hbar),\end{aligned}\tag{12}$$

where i=I, II, III (see Fig.1), $E_{z0}^{\pm}+E_{\parallel}^{\pm}=E_0^{\pm}$, $E_{\parallel}^{\pm}=\frac{\hbar^2k_{\parallel}^2}{2\mu_1}$, $k_{zn}^{\pm}=\sqrt{\frac{2\mu_1}{\hbar^2}(E_{z0}^{\pm}+n\hbar\omega)}$, $c_n^{(i)\pm}$ and $d_n^{(i)\pm}$ are the probability amplitudes in the ith region. The incoming and the outgoing waves (channels) of Floquet scattering form the sidebands (or Floquet channels) with an energy spacing $\hbar\omega$ according to $E_n^{\pm}=E_{\mathrm{F}}^{\pm}+n\hbar\omega$ (n is the sideband index). The mode of $E_n<0$ is an evanescent mode and corresponding sideband is called an evanescent sideband because such a mode with an imaginary k_n cannot propagate.[31,32]

The scattering matrix can be constructed for the time-periodic potential according to Ref.[31]. Similarly, using the continuity of $\Phi_{\pm}$ and $\frac{1}{\mu}\frac{\partial}{\partial z}\Phi_{\pm}$ at the interfaces $z=0$, $z=a_{\mathrm{L}}$,$z=a_{\mathrm{L}}+b$ and $z=a_{\mathrm{L}}+b+a_{\mathrm{R}}$ for Eqs.(11) and (12), the scattering matrix for the double potential wells of Fig.1 is given by

$$S_{\mathrm{total}}=\begin{pmatrix}R & T'\\ T & R'\end{pmatrix},\tag{13}$$

where R and T are matrices whose elements denote the probability amplitudes of reflection and transmission, respectively, for modes incident from the left (region I). For instance, the matrix element t_{nm} is the probability amplitude of transmission for an electron incident from channel m in the region I to appear in channel n in the region III. R' and T' describe the process of scattering incident from the right (region III). Taking $c_n^{(\mathrm{I})\pm}=\delta_{n0}$ and $d_n^{(\mathrm{III})\pm}=0$, the total electron-transmission probabilities of spin "+" and spin "–" components are given by

$$T^{\pm}=\sum_{n=0}^{+\infty}\frac{k_{zn}^{\pm}}{k_{z0}^{\pm}}\left|t_{n0}^{\pm}\right|^2,\tag{14}$$

from which the electron conductance through the double-well structure can be obtained by using the Landauer–Buttiker formula[33,34]

$$G^{\pm}=\frac{e^2}{h}T^{\pm}=\frac{e^2}{h}\sum_{n=0}^{+\infty}\frac{k_{zn}^{\pm}}{k_{z0}^{\pm}}\left|t_{n0}^{\pm}\right|^2.\tag{15}$$

3. Numerical results

We now study numerically the scattering of an incident wave through oscillating double quantum wells and calculate the conductivity. The minimum number of sidebands in the sum of Eq.(14) and Eq.(15) needing to be included may be determined by the oscillation amplitudes of double wells. In general, it is enough to take $N>V_1/(\hbar\omega)$.[35].

The interaction of electrons with the oscillating field leads to a photon-mediated resonance transmission. The incident electrons in the incident channel can drop to the bound states of double wells by means of photon emission. Similarly, the electrons in bound states can also absorb photons and jump to the incident channel or other Floquet channel. This forms the discrete channel of scattering required by the Fano-type resonance. The resonance transmission takes place when the difference in energy between the incident electrons and the bound states of double wells is equal to the energy of one photon or multi-photon. Since the resonance tunnelling through a thin barrier between two wells leads to the level splitting of bound state, the Fano-type resonance spectra through double wells are different from those in the case of single well and are divided into two sets each with two asymmetric resonance peaks due to the Dresselhaus spin–orbit coupling. The difference in phase between two oscillating fields may influence the line shape of transmission

resonance via the asymmetry parameter q. In Fig.(2), we plot the conductivity $G^{\pm}$ as a function of the incident electron energy E_{z0} for $a_{\rm L} = a_{\rm R}$=6 nm , b=15 nm, $V_{\rm L} = V_{\rm R}$= 300 meV, $k_{\|}$=10^6 cm^{-1}, $\hbar\omega$=10 meV, $V_{1\rm L}$=5 meV, $V_{1\rm R}$=5 and 10 meV. The parameters of the semiconductor heterostructure are chosen as γ=0.187eV·nm^3, μ=0.041$m_{\rm e}$ ($m_{\rm e}$ is the mass of a free electron) for GaSb, $\gamma_1 = 8\times10^{-3}$eV·nm^3, μ_1 = 0.081$m_{\rm e}$ for InP according to Ref.[12]. In actual calculation we take $\gamma_1 \approx 0$ because $\gamma \gg \gamma_1$. Sidebands of $n = 0, \pm1, \cdots, \pm10$ are taken into account so that $T^{\pm} = \sum_{n=0}^{10} \frac{k_{zn}^{\pm}}{k_{z0}^{\pm}}|t_{n0}^{\pm}|^2$. The conductivity pattern in Fig.2 shows that double-peak resonances (a_-, b_-) occur approximately at E_z^- = 1.10, 4.57 meV for the spin state "–" and (a_+, b_+) occur approximately at E_z^+ = 3.89, 7.55 meV for the spin state "+". Corresponding bound state energies are E_b^- = −8.90, –5.43 meV and E_b^+ = −6.11, –2.45 meV in the absence of the oscillating fields, and the resonance energies satisfy the relation $E_z^{\pm} - E_b^{\pm} = \hbar\omega$. We can see clearly from Figs.2(a) and 2(b) that the dips of the resonances (a_-, a_+) (corresponding to the low level after splitting) lie on the right-hand side of the asymmetric resonant peaks for the case $\theta_{\rm R} = 0$, while the dips of the resonances (b_-, b_+) (corresponding to the high level after splitting) lie on the left-hand side of the asymmetric resonant peaks. The case of $\theta_{\rm R} = \pi$ (Fig.2(c, d)) is just opposite: the dips of the resonances (a_-, a_+) lie on the left-hand side and the dips of the resonance (b_-, b_+) lie on the right-hand side. The oscillation amplitude $V_{1\rm R}$ not only affects the resonance width Γ, but also suppresses the resonance peaks selectively. When $\theta_{\rm R} = 0$, the resonance peaks (a_-, a_+) may be intensively suppressed by increasing the oscillation amplitude (Fig.2(b)) while the suppressed resonance peaks become (b_-, b_+) (Fig.2(d)) for $\theta_{\rm R} = \pi$.

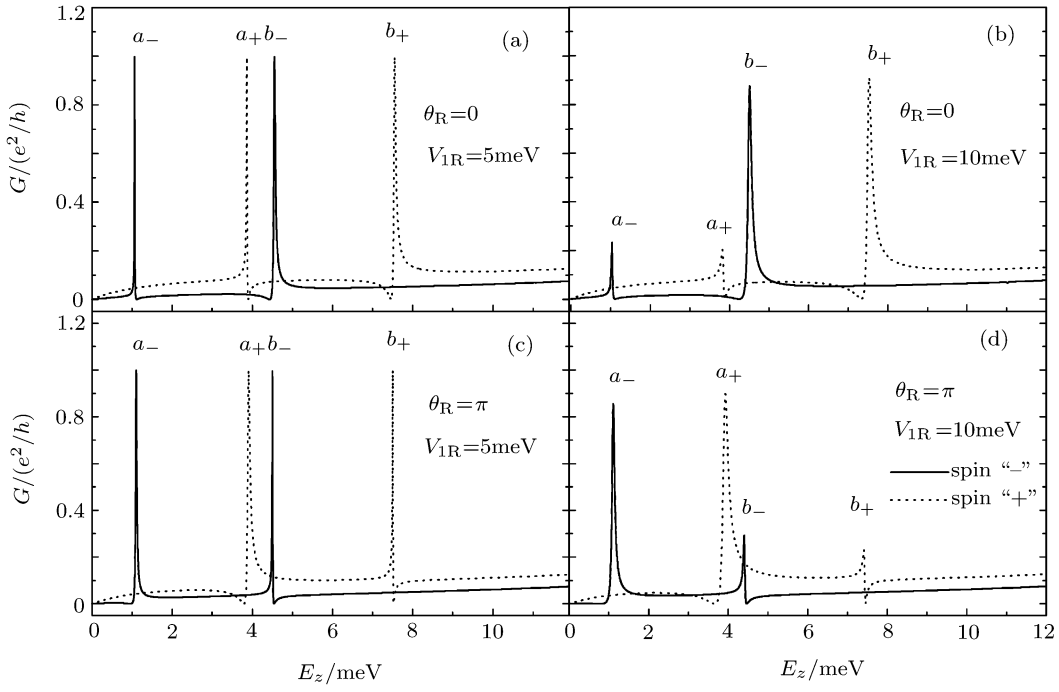


Fig.2. Splitting of Fano-type resonance in conductivity $G^{\pm}$ for $a_{\rm L} = a_{\rm R}$=6 nm, b=15 nm, $k_{\|} = 10^6$ cm^{-1}, $\hbar\omega$=10 meV, $\mu = 0.041m_{\rm e}$, $\mu_1 = 0.081m_{\rm e}$, $\gamma = 0.187eV\cdot$nm^3, $\gamma_1 = 0$, $V_{\rm L} = V_{\rm R}$=300 meV, $\theta_{\rm L} = 0$, $\theta_{\rm R} = 0$ and π, $V_{1\rm L}$=5 meV, $V_{1\rm R}$ = 5 and 10 meV.

Conventionally the asymmetry parameter q is considered to be real, but it is valid only when the system has the time-reversal symmetry (TRS).[25] In our model the TRS is broken due to the difference in phase between two oscillating fields, so Eq.(1) in Fano formula may be expressed as[25,26,36] $G(\epsilon) = G_{\rm d}\frac{(\epsilon + {\rm Re}q)^2 + ({\rm Im}q)^2}{\epsilon^2 + 1}$. When $|{\rm Re}q| \gg |{\rm Im}q|$, $G(\epsilon)$ shows the asymmetric line shapes: a dip lies on the left-hand side of the asymmetric resonant peak for ${\rm Re}q > 0$ and on the right-hand side for ${\rm Im}q < 0$. When $|{\rm Re}q| \ll |{\rm Im}q|$, $G(\epsilon)$ shows the symmetric Breit–

Wigner line shape. In Fig.(3), we plot the phase dependence of conductivity for spin states "+" and "–", where we take $\theta_{\rm L} = 0$ and thus $\theta_{\rm R}$ denotes the difference in phase between two oscillating fields. The asymmetry of resonance peaks shows an interesting characteristic with the variation of $\theta_{\rm R}$. When $\theta_{\rm R}$ increases from 0 to π, the dips of resonances (a_-, a_+) shift from the right-hand side of the asymmetric resonance peaks to the left-hand side. At a critical value of the phase difference value $(\theta_{\rm R}^{\rm c})$ of about $3\pi/10$, the transmission resonance degenerates to a symmetric Breit–Wigner line shape. For the resonances (b_-, b_+), the way is reversed: the dips of resonances shift from the left-hand side to the right-hand side and the symmetric Breit–Wigner resonances occur at a critical value of the phase difference $(\theta_{\rm R}^{\rm c})$ of about $7\pi/10$.

The locations of the asymmetric resonance peaks in energy windows depend not only on the width and depth of the wells, but also on the oscillation frequency of external field. Figure 4 shows the frequency dependence of conductivity for spin states "+" and "–". The asymmetric resonance peaks move toward the direction of higher energy as the frequency increases, which is similar to the case of single well except for the double-peak structure of the transmission spectrum due to the level splitting.

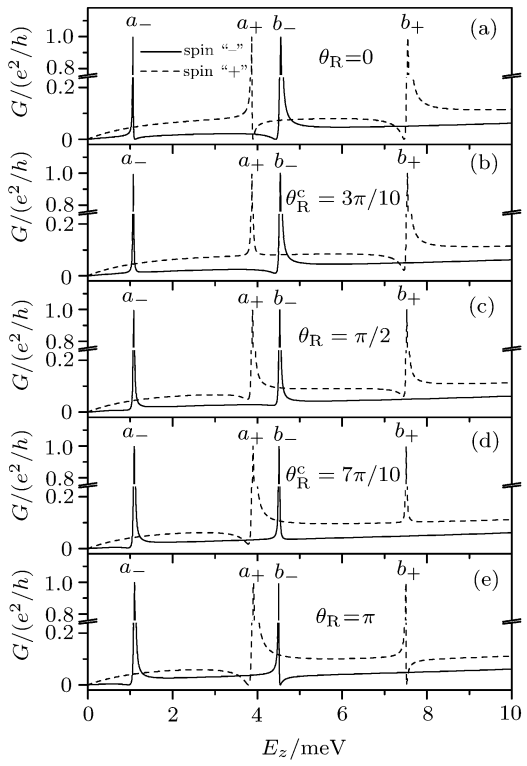


Fig.3. Variation of the Fano-resonance line-shape in conductivity $G^{\pm}$ with $\theta_{\rm R}$.

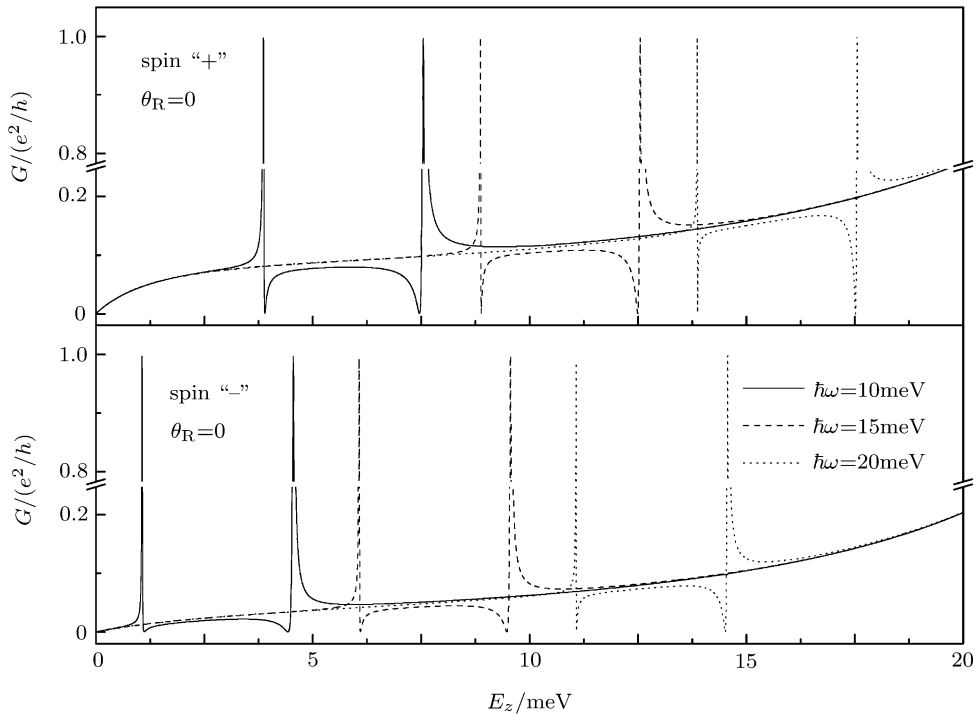


Fig.4. Conductivity $G^{\pm}$ as a function of E_z for different values of $\hbar\omega$ with $\theta_{\rm L} = \theta_{\rm R} = 0$ and $V_{1\rm R}$=5 meV.

The spin-polarization-dependent splitting of Fano-type resonance peak is very useful for realizing the spin polarization current. Figure (5) shows the spin-polarization dependent transmission of electrons as a function of the electron energy according to $P = (T^+ - T^-)/(T^+ + T^-)$. For an energy range ΔE_z from 4 to 6 meV, when $\theta_{\rm R} = 0$, the transmission current is of the spin "–" polarization with a polarization efficiency of some 95%; when $\theta_{\rm R} = \pi$, the transmission current is of the spin "+" polarization with a polarization efficiency of some 75%. There is a contrary property for an energy range ΔE_z from 1.4 to 3.4 meV, the spin "+" polarization is dominant for $\theta_{\rm R} = 0$ and the spin "–" polarization is dominant for

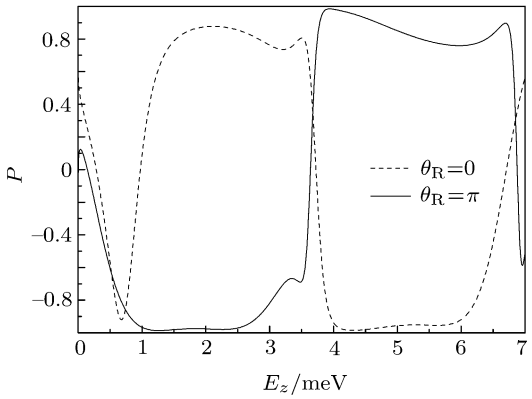


Fig.5. Dependences of polarization efficiency of the transmission electrons on energy of incident electron for $\theta_{\rm R}$=0 and π, with $V_{1\rm L} = V_{1\rm R}$=20 meV.

$\theta_{\rm R} = \pi$. Within a given energy range, the spin polarization of transmission current is completely governed by the difference in phase between two oscillating fields, which suggests that one can exploit the Fano-type resonances of transmission current as the basis of a spin filter and obtain spin polarization current by appropriately matching the energy E_z with the phase difference $\theta_{\rm R}$.

4. Conclusions

We have investigated theoretically the property of the multi-field driven electron transport through double quantum wells in semiconductor heterostructures with spin–orbit coupling. The numerical results demonstrate that the transmission spectra of electrons are divided into two sets due to the level splitting of bound state in double wells and each set contains two asymmetric resonance peaks. When the oscillating amplitude of applied fields, $V_{1\rm R}$, increases, the resonances (a_-, a_+) and (b_-, b_+) may be selectively suppressed by changing the difference in phase between 0 and π. The transmission resonances show the completely different asymmetries for $\theta_{\rm R} = 0$ and π. When the phase difference changes from 0 to π, the dips of asymmetry resonances move from one side of resonance peaks to the other side and at a critical value of phase difference $\theta_{\rm R}^{\rm c}$ the asymmetric Fano resonance degenerates into the symmetric Breit–Wigner resonance. Within a given energy range, the spin polarization of transmission current is completely governed by the difference in phase between the two oscillating fields and thus an appropriate match of the energy E_z with the phase difference $\theta_{\rm R}$ may be used to realize the tunable spin filtering.

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