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Electron resonance-transmission through a driven quantum well with spin–orbit coupling*

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We have studied the spin-dependent electron transmission through a quantum well driven by both dipole-type and homogeneous oscillating fields. The numerical evaluations show that Dresselhaus spin–orbit coupling induces the splitting of asymmetric Fano-type resonance peaks in the conductivity, in which the dipole modulation and the homogeneous modulation are equivalent. Therefore, we predict that the dipole-type oscillation, which is more practical in the experimental setup, can be used to realize the tunable spin filters by adjusting the field oscillation-frequency and the amplitude as well.

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

1. Introduction

Quantum transport in driven mesoscopic systems has attracted much interest because of a variety of novel phenomena and effects.[1,2] The photon-assisted-tunnelling (PAT) effect without considering the spin degree of freedom of electrons has been observed in quantum superlattices,[3] tunnelling diodes,[4] and split gate in quantum dots.[5] Recently, a large number of spin-dependent physical phenomena, for instance, Datta–Das spin field-effect transistor,[6] spin transport,[7] spin Hall effect[8] and so on, have attracted great attention due to the prospects for technological applications. Spin polarized transport of electrons in periodically driven mesoscopic systems also displays many interesting phenomena, such as spin-dependent Fano-type resonance and tunable spin-filtering.[9] Fano-type resonance, which may trace back to the autoionization phenomenon of atoms,[10] arises from interference between the bound states and the continuum of band. The first tunable Fano experiment[11] is realized in an Aharonov–Bohm ring with a quantum dot embedded in one of its arms. The quantum transport in periodically driven semiconductor heterostructure exhibits fascinating Fano-type resonance which splits into two asymmetric resonance peaks due to Dresselhaus spin–orbit interaction.[9] The time-periodic field mixes the continuum and bound states, and leads to the appearance of asymmetric Fano-type resonance. While the spin–orbit coupling induces spin-dependent splitting in resonance transmission, which plays an important role in creating spin-polarization-dependent transmission electron-current.

As is well known, a more practical oscillation-type is dipole-oscillation in most experimental setup, and the dipole-type is easier to exploit experimentally than the spatially uniform field. In this paper, we investigate the spin-dependent electron transmission through a semiconductor quantum well with a dipole-type oscillation-field and a homogeneous oscillation-field. The quantum well in InP–GaSb–InP semiconductor heterostructures formed by a thin small band gap layer with large Dresselhaus spin–orbit coupling sandwiched between two larger band gap layers with very small spin–orbit coupling. Electrons pass through three regions from the left to the right, I, II, and III, with I corresponding to the left region, III to the right, and II to the central well region. Our numerical results show that: (1) each of the two modulation leads to

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Fano-type resonance in the transmission spectrum of electrons, (2) spin–orbit coupling induces Fano-type resonance splitting, and (3) the dipole modulation and the homogeneous modulation are equivalent under a certain condition.

In Section 2, we introduce our model system and solve the time-periodic Schrödinger equation using the Floquet theorem, and obtain the transmission probability and the spin-dependent conductance. In Section 3, we study numerically the behaviour of electron transmission through a well with spin–orbit coupling and oscillating field. Finally we make some concluding remarks In Section 4.

2. The model and formalism

We consider the transmission of electron with an incident wave vector $\boldsymbol{k} = \left(k_{||}\cos\varphi, k_{||}\sin\varphi, k_z\right)$ above a one-dimension all time-periodic potential well which extends from 0 to a. The potential is

$$V(z,t) = \begin{cases} 0, & z < 0 \text{ and } z > a; \\ -V_0 + V_1\cos(\omega t) + eF_1 z\cos(\omega t), & 0 \le z \le a. \end{cases} \tag{1}$$

Here φ is the polar angle of the wave vector $\boldsymbol{k}$ in the XY plane, $k_{||}$ is the wave vector in the plane of the well, and k_z is the wave vector normal to the well pointing along the direction of tunnelling; well depth V_0 grows along $z \parallel [001]$, F_1 is a uniform oscillating electric field and V_1 is a homogeneous potential modulation. Considering a single electron model, we can neglect the electron–electron interactions and assume the temperature is low enough that electron–phonon interactions[12] can be neglected. Thus the electron motion in semiconductor heterostructure may be described by a time-dependent Schrödinger equation

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\phi(z,t) = \hat{H}\phi(z,t) \tag{2}$$

with Hamiltonian

$$\hat{H} = -\frac{\hbar^2}{2\mu}\frac{\partial^2}{\partial z^2} + \frac{\hbar^2 k_{||}^2}{2\mu} + V(z,t) + \hat{H}_{\mathrm{D}}. \tag{3}$$

We assume the kinetic energy of the incident electron to be much smaller than the well depth V_0, then the Dresselhaus term $\hat{H}_{\mathrm{D}}$ may be simplified to[13]

$$\hat{H}_D = \begin{cases} 0, & z < 0 \text{ and } z > a; \\ \gamma\left(\hat{\sigma}_x k_x - \hat{\sigma}_y k_y\right)\dfrac{\partial^2}{\partial z^2}, & 0 \le z \le a, \end{cases} \tag{4}$$

where μ is the electron effective mass, $\hat{H}_{\mathrm{D}}$ the spin-dependent k^3 Dresselhaus term in Zinc-blende structure semiconductors,[14] $\hat{\sigma}_{\mathrm{i}}$ the Pauli matrices, and γ a material constant describing the strength of the spin–orbit coupling. We have assumed $\gamma \approx 0$ outside the well. The Dresselhaus term (4) is diagonalized by the spinors

$$\chi_s = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 \\ -s\mathrm{e}^{-\mathrm{i}\varphi} \end{pmatrix}, \tag{5}$$

which describes the electron spin states with $s = \pm 1$ corresponding to the opposite spin directions. In the spin subspace the Hamiltonian $\hat{H}$ may be reduced as

$$\hat{H}_s = \chi_s^{+}\hat{H}\chi_s = -\frac{\hbar^2}{2\mu_s}\frac{\partial^2}{\partial z^2} + \frac{\hbar^2 k_{||}^2}{2\mu} + V(z,t), \tag{6}$$

where the modified effective mass of electron depends on Dresselhauss coupling constant γ, the in-plane electron wave vector $k_{||}$ and the orientation of the electron spin, and is given by $\mu_s = \mu\left(1 + s\gamma 2\mu k_{\|}/\hbar^2\right)^{-1}$. Although the modification for the electron effective mass is very small, it plays an important role in inducing the splitting of asymmetric Fano resonance peak. Thus the Schrödinger equation can be rewritten as

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\phi_s(z,t) = \hat{H}_s\phi_s(z,t) \tag{7}$$

with

$$\phi_s(z,t) = \chi_s\psi_s(z,t)\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}), \tag{8}$$

where $\boldsymbol{\rho} = (x,y)$ is a vector in the plane of the well. Inserting Eqs.(6) and (8) into Eq.(7) we obtain the equation for $\psi_s(z,t)$

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\psi_s(z,t) = -\frac{\hbar^2}{2\mu_s}\frac{\partial^2}{\partial z^2}\psi_s(z,t) + \left[\frac{\hbar^2 k_{\|}^2}{2\mu} + V(z,t)\right]\psi_s(z,t). \tag{9}$$

Inside the potential well, the potential $V(z,t)$ is a periodic function of time, according to the Floquet

theorem,[15] Eq.(9) has a solution of the form

$$\psi_s(z,t) = \Psi_s(z,t)\exp(-\mathrm{i}E_F^s t/\hbar), \tag{10}$$

where $E_{\mathrm{F}}^{\mathrm{s}}$ is the Floquet energy and $\Psi_s(z,t)$ is a periodic function of time: $\Psi_s(z,t) = \Psi_s(z,t+T)$ with period $T = 2\pi/\omega$. A Floquet state is the analogue to a Bloch state when replacing a spatially periodic potential with a time periodic one. The quasimomentum of the Bloch state becomes the quasienergy $E_{\mathrm{F}}^{\mathrm{s}}$ of the Floquet state. In order to solve Eq.(9), we write the potential $V(z,t) = g(t) + zf(t)$, with $g(t) = -V_0 + V_1\cos(\omega t)$ and $f(t) = eF_1\cos(\omega t)$, and define $p(t) = \int^t \mathrm{d}\tau f(\tau)$. According to Refs.[1, 16] the wavefunction $\Psi_s(z,t)$ can be expressed as

$$\Psi_s(z,t) = \Phi_s(\xi,t)\eta_s(z,t), \tag{11}$$

where

$$\xi = z - \frac{eF_1}{\mu_s\omega^2}\cos(\omega t), \tag{12}$$

$$\eta_s(z,t) = \exp\left\{-\frac{\mathrm{i}}{\hbar}\left(E - E_{\mathrm{F}}^{\mathrm{s}} - \frac{e^2F_1^2}{4\mu_s\omega^2}\right)t - \frac{\mathrm{i}eF_1\sin(\omega t)}{\hbar\omega} + \frac{\mathrm{i}e^2F_1^2\sin(2\omega t)}{8\hbar\mu_s\omega^3}\right\}, \tag{13}$$

with E being energy constant. Substituting Eqs.(11) and (13) into Eq.(9), we obtain the following equation for $\Phi_s(\xi,t)$:

$$\mathrm{i}\hbar\frac{\partial}{\partial t}\Phi_s(\xi,t) = \left[-\frac{\hbar^2}{2\mu_s}\frac{\partial^2}{\partial\xi^2} + \frac{\hbar^2k_{\parallel}^2}{2\mu}\right]\Phi_s(\xi,t) + \left[-V_0 - E_{\mathrm{F}}^{\mathrm{s}} + V_1\cos(\omega t) + \frac{e^2F_1^2}{4\mu_s\omega^2}\right]\Phi_s(\xi,t). \tag{14}$$

Let

$$\Phi_s(\xi,t) = \varphi_s(\xi)\exp\left[-\frac{\mathrm{i}}{\hbar\omega}V_1\sin(\omega t)\right]. \tag{15}$$

Substitute Eq.(15) into Eq.(14), $\varphi_s(\xi)$ satisfies a simple stationary state Schrödinger equation:

$$\left[-\frac{\hbar^2}{2\mu_s}\frac{\partial^2}{\partial\xi^2} + \frac{\hbar^2k_{\parallel}^2}{2\mu} - V_0\right]\varphi_s(\xi) = E\varphi_s(\xi). \tag{16}$$

Thus the wavefunction inside the oscillation region II can be expressed as:

$$\begin{aligned}\phi_s^{\mathrm{II}}(z,t) =& \chi_s\mathrm{e}^{-\frac{\mathrm{i}}{\hbar}E_{\mathrm{F}}^{\mathrm{s}}t}\sum_{l=-\infty}^{+\infty}[a_l^s\mathrm{e}^{\mathrm{i}q_l^s\xi} + b_l^s\mathrm{e}^{-\mathrm{i}q_l^s\xi}]\\ &\times\exp\left[-\frac{\mathrm{i}}{\hbar}l\hbar\omega t + \frac{\mathrm{i}e^2F_1^2\sin(2\omega t)}{8\hbar\mu_s\omega^3}\right.\\ &\left.-\mathrm{i}\left(\frac{V_1}{\hbar\omega} + \frac{eF_1z}{\hbar\omega}\right)\sin(\omega t)\right]\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}),\end{aligned} \tag{17}$$

where a_l^s and b_l^s are the two scattering coefficients to be determined. The sideband wave vectors q_l^s in solution (17) satisfy the relation $q_l^s = \sqrt{\frac{2\mu_s}{\hbar^2}\left(E_{\mathrm{F}}^{\mathrm{s}} - \frac{e^2F_1^2}{4\mu_s\omega^2} + l\hbar\omega - \frac{\hbar^2k_{\parallel}^2}{2\mu} + V_0\right)}$. Substituting Eq.(12) into Eq.(17) and using the following identities:

$$\begin{aligned}\exp[-\mathrm{i}z\sin(\omega t)] =& \sum_m J_m(z)\exp(-\mathrm{i}m\omega t),\\ \exp[-\mathrm{i}z\cos(\omega t)] =& \sum_m(-i)^m J_m(z)\exp(-\mathrm{i}m\omega t),\end{aligned} \tag{18}$$

with $J_m(z)$ the Bessel function of the first kind, Eq.(17) can be written as

$$\begin{aligned}\phi_s^{\mathrm{II}}(z,t) =& \chi_s\sum_{n=-\infty}^{+\infty}\sum_{l=-\infty}^{+\infty}[a_l^s\,\Phi_{snl}^{+}(z) + b_l^s\,\Phi_{snl}^{-}(z)]\\ &\times\exp\left[-\frac{\mathrm{i}}{\hbar}E_{zn}^s t\right]\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho} - \mathrm{i}E_{\parallel}^s t/\hbar)\end{aligned} \tag{19}$$

with

$$\Phi^{\pm}_{snl}(z) = \sum_{\alpha=-\infty}^{+\infty}\sum_{\beta=-\infty}^{+\infty} \mathrm{e}^{\pm \mathrm{i}q^s_l z}(-\mathrm{i})^{\alpha} J_{\alpha}\left(\frac{\pm q^s_l e F_1}{\mu_s \omega^2}\right) J_{\beta}\left(\frac{e^2F_1^2}{8\hbar\mu_s\omega^3}\right) \times J_{n-l-\alpha+2\beta}\left(\frac{V_1}{\hbar\omega}+\frac{eF_1 z}{\hbar\omega}\right). \tag{20}$$

Since electrons incident to the oscillating region will be scattered inelastically into an infinite number of Floquet sidebands with energy spacing of $\hbar\omega$, the wavefunction outside the potential can be written as the wave superpositions of all kinds of energy:

$$\phi^{\mathrm{I}}_s(z,t) = \chi_s\left[\mathrm{e}^{\mathrm{i}k^s_{z0}z-\frac{\mathrm{i}}{\hbar}E^s_{z0}t} + \sum_{n=-\infty}^{+\infty} r^s_{n0}\mathrm{e}^{-\mathrm{i}k^s_{zn}z-\mathrm{i}E^s_{zn}t/\hbar}\right]\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}-\mathrm{i}E^s_{\|}t/\hbar), \tag{21}$$

$$\phi^{\mathrm{III}}_s(z,t) = \chi_s\sum_{n=-\infty}^{+\infty} t^s_{n0}\mathrm{e}^{\mathrm{i}k^s_{zn}z-\mathrm{i}E^s_{zn}t/\hbar}\exp(\mathrm{i}\boldsymbol{k}_{||}\cdot\boldsymbol{\rho}-\mathrm{i}E^s_{\|}t/\hbar) \tag{22}$$

with $E^s_{z0}+E^s_{\|}=E^s_0$, $E^s_{\|}=\dfrac{\hbar^2k^2_{\|}}{2\mu_1}$ and $k^s_{zn}=\sqrt{\frac{2\mu_1}{\hbar^2}(E^s_{z0}+n\hbar\omega)}$. For simplicity, we restrict our discussion to $E^s_0<\hbar\omega$ which is the Floquet energy of the propagating mode with the lowest energy. r^s_{n0} and t^s_{n0} are the probability amplitudes of the reflecting waves and outgoing waves from the sideband 0 to sideband n, respectively. The incoming and outgoing waves (channels) of Floquet scattering form the sidebands (or Floquet channels) with energy spacing $\hbar\omega$ according to $E^s_n=E^{\mathrm{s}}_{\mathrm{F}}+n\hbar\omega$ (n is the sideband index). The mode of $E_n<0$ is an evanescent mode and the corresponding sideband is called an evanescent sideband because such a mode with imaginary k_n cannot propagate.[17] The continuity of ϕ_s and $\dfrac{1}{\mu}\dfrac{\partial}{\partial z}\phi_s$ at the interfaces $z=0$ and $z=a$ requires

$$\begin{aligned}
&\sum_{l=-\infty}^{+\infty}([\Phi^{+}_{snl}(0)a^s_l+\Phi^{-}_{snl}(0)b^s_l]=\delta_{n0}+r^s_{n0},\\
&\frac{\mu_1}{\mu_s}\sum_{l=-\infty}^{+\infty}[\bar{\Phi}^{+}_{snl}(0)a^s_l+\bar{\Phi}^{-}_{snl}(0)b^s_l]=\mathrm{i}k^s_{zn}(\delta_{n0}-r^s_{n0}),\\
&\sum_{l=-\infty}^{+\infty}([\Phi^{+}_{snl}(a)a^s_l+\Phi^{-}_{snl}(a)b^s_l]=\mathrm{e}^{\mathrm{i}k^s_n a}t^s_{n0},\\
&\frac{\mu_1}{\mu_s}\sum_{l=-\infty}^{+\infty}[\bar{\Phi}^{+}_{snl}(a)a^s_l+\bar{\Phi}^{-}_{snl}(a)b^s_l]=\mathrm{i}k^s_{zn}\mathrm{e}^{\mathrm{i}k^s_{zn}a}t^s_{n0}
\end{aligned} \tag{23}$$

with $\bar{\Phi}^{\pm}_{snl}(0)=\dfrac{\partial\Phi^{\pm}_{snl}}{\partial z}|_{z=0}$ and $\bar{\Phi}^{\pm}_{snl}(a)=\dfrac{\partial\Phi^{\pm}_{snl}}{\partial z}|_{z=a}$. The continuity conditions of wavefunctions Eq.(23) can be expressed in matrix forms

$$\begin{aligned}
&M^+_sA^s+M^-_sB^s=\Delta+R^s,\\
&\frac{\mu_1}{\mu_s}(\breve{M}^+_sA^s+\breve{M}^-_sB^s)=K_s(\Delta-R^s),\\
&N^+_sA^s+N^-_sB^s=S_sT^s,\\
&\frac{\mu_1}{\mu_s}(\breve{N}^+_sA^s+\breve{N}^-_sB^s)=K_sS_sT^s
\end{aligned} \tag{24}$$

with the help of the square matrices defined by the matrix elements $M^{\pm}_{snl}=\Phi^{\pm}_{snl}(0)$, $\breve{M}^{\pm}_{snl}=\bar{\Phi}^{\pm}_{snl}(0)$, $N^{\pm}_{snl}=\Phi^{\pm}_{snl}(a)$, $\breve{N}^{\pm}_{snl}=\bar{\Phi}^{\pm}_{snl}(a)$, $K_{snl}=\mathrm{i}k^s_{zl}\delta_{nl}$, and $S_{snl}=\mathrm{e}^{\mathrm{i}k^s_l a}\delta_{nl}$ and the column matrices by the matrix elements $A^s_l=a^s_l$, $B^s_l=b^s_l$, $\Delta_l=\delta_{l0}$, $R^s_l=r^s_{l0}$, and $T^s_l=t^s_{l0}$. Here R^s and T^s are the matrices of the reflection and the transmission amplitude. From the matrix equation (24) one can obtain the matrix of the transmission amplitude:

$$T^s=2S^{-1}_sN_sK_s\Delta, \tag{25}$$

where $N_s=\left[N^-_s-N^+_s(P^+_s)^{-1}P^-_s\right](D_s)^{-1}\ (Q^+_s)^{-1}$ with

$$\begin{aligned}
P^{\pm}_s=&K_sN^{\pm}_s-\frac{\mu_1}{\mu_s}\breve{N}^{\pm}_s,\\
Q^{\pm}_s=&K_sM^{\pm}_s+\frac{\mu_1}{\mu_s}\breve{M}^{\pm}_s,\\
D_s=&(Q^+_s)^{-1}Q^-_s-(P^+_s)^{-1}P^-_s.
\end{aligned}$$

The total electron-transmission probabilities of the opposite spin components are given by

$$T^s=\sum_{n=0}^{+\infty}\frac{k^s_{zn}}{k^s_{z0}}\left|t^s_{n0}\right|^2, \tag{26}$$

from which the conductance of the electrons through semiconductor heterostructures can be obtained by Landauer–Buttiker formula[18]

$$G^s = \frac{e^2}{h}T^s = \frac{e^2}{h}\sum_{n=0}^{+\infty}\frac{k_{zn}^s}{k_{z0}^s}\left|t_{n0}^s\right|^2. \tag{27}$$

The spin polarization efficiency of the transmitted electron can be found from

$$P = \frac{T^+ - T^-}{T^+ + T^-}. \tag{28}$$

3. Numerical results

We now study numerically the scattering of an incident wave through an oscillating quantum well with spin–orbit coupling and calculate the conductivity. The interaction of electrons with the oscillating field leads to photon-mediated transmission resonances. The incident electron can drop to the bound states in the well by means of photon emission. Similarly, the electron in bound states can also absorb photons and jump to the incident channels or other Floquet channels. Transmission resonance takes place when the energy difference between the incident electrons and the bound states is equal to the energy of one photon or more photons. Since the number of equations to determine the coefficients a_l^s, b_l^s, r_{n0}^s and t_{n0}^s is infinite, we have to make choice of the effective upper bound in order to accomplish practical calculations. The minimum number of sidebands needed to be included is determined by the oscillation amplitude of the potential well. In general, it is enough to take $N > \max\left(\frac{q_0^s eF_1}{\mu_s\omega^2}, \frac{e^2F_1}{8\hbar\mu_s\omega^3}, \frac{V_1}{\hbar\omega} + \frac{eF_1a}{\hbar\omega}\right)$, and when $F_1 = 0$, to take $N > V_1/(\hbar\omega)$.[19]

In Fig.1, we plot the conductivity $G^{\pm}$ as a function of incident electron energy E_{z0} for V_0=300meV, a=39.6 nm, $k_{\|} = 10^6\text{cm}^{-1}$, $\hbar\omega$=10meV, $V_1 = 0$ and F_1=0.1, 0.2, 0.4mV/nm. The parameters of the semiconductor heterostructure are chosen as γ=0.187eV·nm^3, μ=0.041m_e (m_e being the mass of the free electron) for GaSb, $\gamma_1 = 8\times10^{-3}$eV·nm^{-3}, $\mu_1 = 0.081m_e$ for InP they are taken from Ref.[14]. In actual calculation we take $\gamma_1 \approx 0$ since $\gamma \gg \gamma_1$. Sidebands of $n = 0, \pm1, \cdots, \pm10$ are taken into account in $T^{\pm} = \sum_{n=0}^{20}\frac{k_{zn}^{\pm}}{k_{z0}^{\pm}}|t_{n0}^{\pm}|^2$. The conductivity pattern in Fig.1 shows that the two asymmetric resonance peaks appear approximately at E_z^+=9.05 meV for the spin-up state + and E_z^-=1.66 meV for the spin-down state –, and the corresponding bound state energies are $E_{\text{b}}^+ = -0.94$meV, $E_{\text{b}}^- = -8.34$meV in the absence of oscillating field, and the resonance energies satisfy the relation $E_z^{\pm} = \hbar\omega + E_{\text{b}}^{\pm}$. The amplitude of dipole oscillating field reflects the strength of electron coupling with the external field. The correspondence between the resonance location and the energy level in the well makes sense only when, strictly speaking, $F_1 \to 0$. The electron coupled with the applied fields leads to the broadening of the level and thus the asymmetric resonance can take place within a certain range of energy rather than at a single level. Figures 1(a)–1(c) shows the variation of the spectrum in which the asymmetric resonance peaks are getting "fat" with increasing dipole oscillation amplitude F_1. In addition, the amplitude of dipole oscillating field also reflects the capability of electron emitting and absorbing photons(exchanging photons). When E_z increases continuously and enters the second incident channel, i.e. $\hbar\omega < E_z < 2\hbar\omega$, Fig.1(c) shows that the asymmetric resonance peaks of two-photon process appears.

The transmission patterns of electrons is shown in Fig.2 under the homogeneous oscillation condition, for V_1=1.98, 3.96, 7.92meV and $F_1 = 0$. The homogeneous oscillation strength is chosen as the average dipole oscillation strength in Fig.1: $V_1 = eF_1a/2$ with F_1=0.1, 0.2, 0.4mV/nm, which results in an interesting result that the dipole modulation and the homogeneous modulation are completely equivalent when $V_1 = eF_1a/2$. In Fig.3, we plot the frequency dependence of the conductivity for the dipole oscillation (see (a) and (b)) or the homogeneous oscillation (see (c) and (d)), from which one again finds that the dipole modulation and the homogeneous modulation are also equivalent for different frequencies if only $V_1 = eF_1a/2$. It is apparent that the control of resonance location is much easier by adjusting the dipole field-frequency than the parameters of the semiconductor heterostructure in practical experiments.

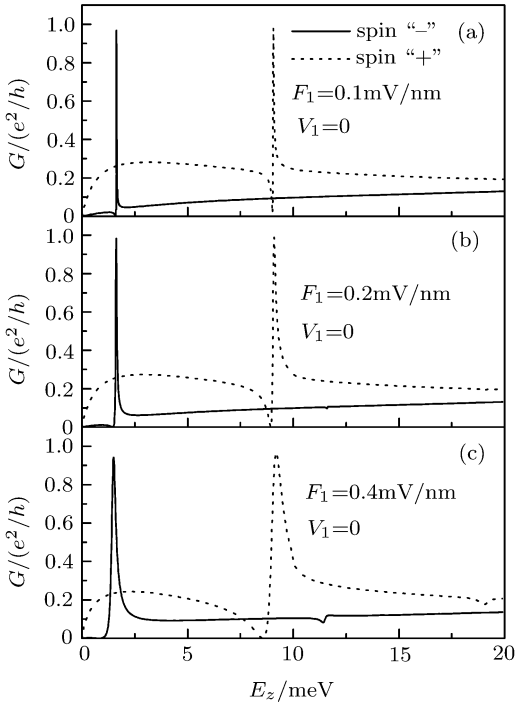


Fig.1. Splitting of Fano-type resonance in conductivity $G^{\pm}$, for a=39.6nm, $k_{\parallel} = 10^6\text{cm}^{-1}$, $\hbar\omega$=10 meV, $\mu = 0.041m_e$, $\mu_1 = 0.081m_e$, γ=0.187eV·nm^3, γ_1=0, V_0=300meV, F_1=0.1, 0.2, 0.4mV/nm, V_1=0.

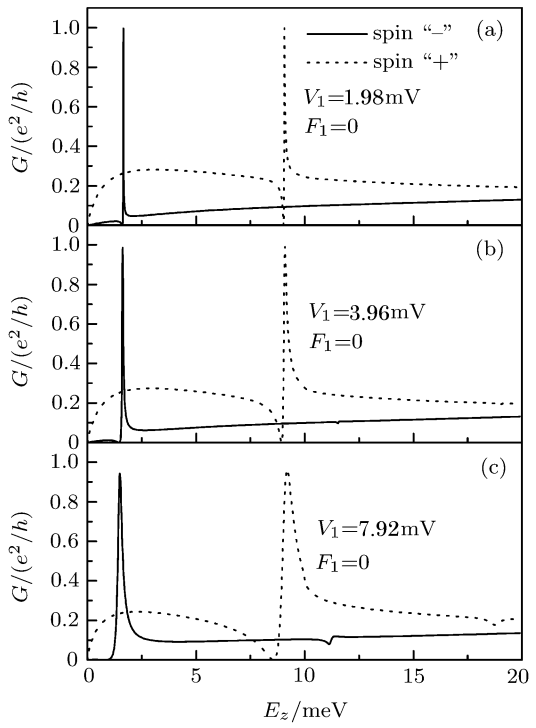


Fig.2. Splitting of Fano-type resonance in conductivity $G^{\pm}$ for $F_1 = 0$, V_1=1.98, 3.96, 7.92meV and the rest parameters are the same as those in Fig.1.

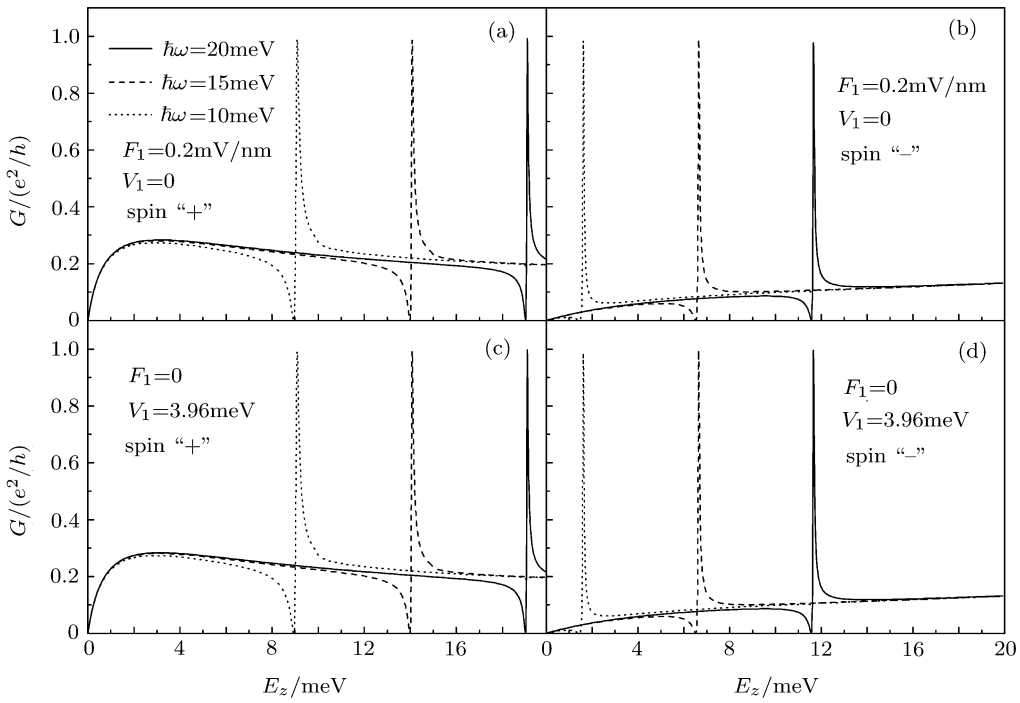


Fig.3. Conductivity $G^{\pm}$ as a function of E_z for different oscillating frequencies for the same parameters as in Fig.1. (a) and (b) for F_1=0.2mV/nm and $V_1 = 0$; (c) and (d) for $F_1 = 0$ and V_1=3.96meV.

Figure 4 shows the spin-polarization efficiency of electrons as a function of the incident electron energy for different oscillating amplitudes. The top panel corresponds to the result obtained with the absence of homogeneous modulation and the bottom panel to that with the absence of dipole modulation. The identical curves in the two panels show once again the dipole modulation and the homogeneous modulation are also equivalent for $V_1 = eF_1a/2$. Moreover we notice that the polarization efficiency of the transmission electrons may reach 90% for the small oscillating amplitude of the applied field and the low incident energy of electron (for example, E_z <5meV, F_1=0.4mV/nm or V_1=7.92meV for $\hbar\omega$=15meV). Thus we expect that the more practical dipole-type oscillation in the experimental setup can be used to realize tunable spin filters by adjusting the oscillation frequency and the amplitude of the dipole external field.

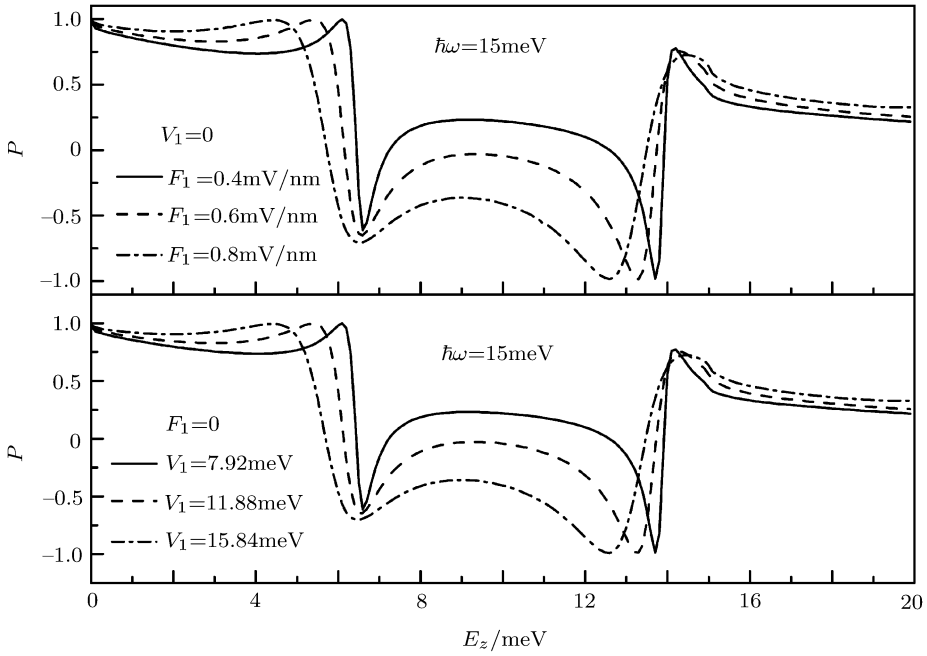


Fig.4. Dependence of polarization efficiency of the transmission electrons on the energy of incident electron for $\omega\hbar$=15meV and the rest parameters are the same as those in Fig.1.

4. Conclusions

In summary, we have investigated theoretically the properties of the spin-dependent electron transport through a quantum well driven by both dipole-type and homogeneous oscillation fields. The numerical results demonstrate that Dresselhaus spin–orbit coupling induces the splitting of asymmetric Fano-type resonance peaks in the conductivity, in which the dipole modulation and the homogeneous modulation are equivalent in the strength-relation $V_1 = eF_1a/2$. Therefore, the dipole-type oscillation-field can be used to realize the tunable spin filters by adjusting the field oscillation-frequency and amplitude.

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