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Spin transmission through a mesoscopic zigzag Rashba wire

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Abstract
The ballistic transport through a mesoscopic zigzag Rashba wire is studied via the Green function technique and the spin transmission is investigated with an unpolarized charge current injected. These results are compared with those of a straight wire, and show that to generate a nonzero spin conductance in the longitudinal direction, the mirror symmetry with respect to that direction must be broken. This necessary condition is satisfied in the zigzag wire. Its ballistic transport properties are also affected by the spin-dependent quantum interference effect, resulting from the Rashba spin–orbit coupling and scattering at corners and at interfaces between the wire and the leads. Mixing between different pairs of transverse modes is crucial for spin transmission, and only if more than one pair of modes are open can a measurable spin-polarized current be obtained in the longitudinal direction. These results provide one way to generate an artificially controllable spin-polarized current in mesoscopic Rashba systems.

1. Introduction

One of the most important possibilities of nanoelectronics is the hope of using spin, in addition to charge, for quantum information processing [1–3]. The basis of this application is the generation of spin-polarized current and quantum control of coherent spin states. In several proposals presented to surmount these two problems, taking advantage of the spin–orbit (SO) coupling in semiconductors [4–7] has attracted a lot of attention. Due to inversion asymmetry of the confining potential for a two-dimensional electronic gas (2DEG) in semiconductor heterostructures, the Rashba SO coupling [8] plays an important role in electronic transport. Although it may be erased by disorder in a bulk system, a pure spin current appears in the transverse direction—the so-called spin Hall effect (SHE)—with an unpolarized electronic current flowing through a mesoscopic Rashba system. This fact has been demonstrated theoretically in mesoscopic Rashba rings [9, 10] and rectangular Rashba planes with their sizes smaller than the coherence length [11–13]. However, in these regular structures, the transparency of a spin-up electron along the longitudinal direction is always equal to that of a

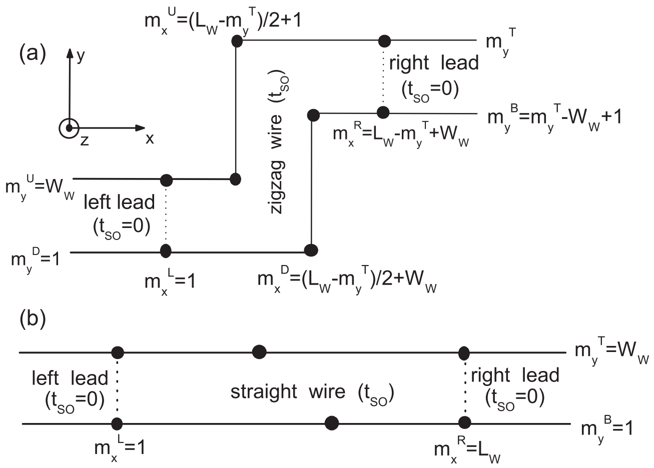

Figure 1. Schematic illustration of the structure.

spin-down electron. The coexistence of the SHE and the zero longitudinal spin conductance of the z-component $\sigma_z = 0$ in the same systems is noticeable.

The Hamiltonian of 2DEG with Rashba SO coupling is [8]

$$H = \frac{\vec{p}^2}{2m^*} + \frac{\alpha}{\hbar}(\vec{\tau} \times \vec{p}) \cdot \vec{e}_z + V_{\rm conf}(x, y). \tag{1}$$

Here, $\vec{\tau}$ is the Pauli operator, α the strength of the Rashba SO coupling and $V_{\rm conf}$ the potential confining electrons to a mesoscopic region within the 2DEG. With the space-inversion symmetry broken, the zero longitudinal spin conductance cannot be derived solely from the time-reversal symmetry. Obviously, it must come from other high structural symmetries possessed by those regular structures. In fact, the Hamiltonians of rings and rectangular planes are invariant under reflection with respect to two mirrors: one is along the longitudinal direction and the other is along the transverse direction. (Their Hamiltonians are also invariant under C_2 rotation.) The former mirror symmetry results in the zero σ_z as we can see from the discussion in section 3. Only structures without this type of mirror symmetry can generate nonzero σ_z. In the present paper, spin transmission through a mesoscopic zigzag Rashba wire (cf figure 1) is studied, where both of the two types of mirror symmetries are destroyed (but the C_2 rotation symmetry is reserved), and that necessary condition is satisfied in this system.

In a straight quantum wire with $\alpha = 0$, a series of spin-degenerate transverse modes can be formed. With the straight wire turned into a zigzag one, different transverse modes can be mixed by corner scattering as well as by interface scattering. With the Rashba SO coupling introduced, the spin degeneracy is lifted and $|\uparrow\rangle$ and $|\downarrow\rangle$ are not the eigenstates of the system. When an electron passes through a 2DEG with Rashba SO coupling, its spin oscillates on the scale of the spin precession length $L_{\rm SO} = \frac{\pi\hbar^2}{2m^*\alpha}$. In ballistic transport through mesoscopic Rashba systems, the spin-dependent quantum interference effect [9, 14–20] plays an important role, and this interference effect can be artificially controlled by adjusting the Rashba SO coupling strength α. The purpose of the present paper is to demonstrate that an artificially controllable spin-polarized current can be generated in the longitudinal direction through a mesoscopic zigzag Rashba wire.

For this purpose, we study the spin transmission through the zigzag wire via the Green function technique [11, 12, 21, 22] together with the formalism of principal layers in the framework of the surface Green function matching theory [21–24], assuming that the structural dimension is much smaller than the coherence length. The results of a straight wire with the same width and length are also obtained as a comparison. Here, for numeric application of the Green function technique, a tight-binding Hamiltonian, rather than the continuum model (1), is adopted. As previous works show [9–13], these two models can give the same results. According to our numeric results, with the mirror symmetries broken, nonzero σ_z can really be found in this zigzag wire. Mixing between different pairs of modes is crucial for spin transmission, and only if more than one pair of modes participate in transport can an experimentally measurable spin conductance be obtained in the longitudinal direction. This nonzero σ_z oscillates with the Rashba SO coupling. Our theoretical results provide one way to generate an artificially controllable spin-polarized current in mesoscopic Rashba systems.

The organization of this paper is as follows. In section 2, the theoretical model and the calculation method are presented. In section 3, the numerical results are illustrated and discussed. A brief summary is given in section 4.

2. Model and formulae

In the present paper, we consider the ballistic transport through a mesoscopic zigzag Rashba wire and investigate the spin transmission with an unpolarized charge current injected through ideal leads. The results for a straight wire with the same length and width as the zigzag one are presented as a comparison to show the influence of structural symmetry. Our attention is also paid to the influence of corner scattering in the zigzag structure and how σ_z varies with the Rashba SO coupling.

The corresponding zigzag and straight wires are schematically illustrated in figure 1. In these two Rashba wires, the SO hopping parameter is $t_{\rm SO}$. Both of these two structures are $L_{\rm W}$ in length and $W_{\rm W}$ in width, and each of them are connected with two ideal leads of the same width $W_{\rm W}$. In these ideal leads, no SO coupling exists. The tunnelling matrix element between the wire and the leads is $t_{\rm L}$. The Hamiltonian of the system can be written in the tight-binding representation as [11–13]

$$H = H_{\rm L} + H_{\rm R} + H_{\rm W} + H_{\rm T}, \tag{2}$$

where $H_{\rm L(R)}$, $H_{\rm W}$ and $H_{\rm T}$ are the Hamiltonians of the left (right) lead, wire and tunnelling between them. They are

$$H_{\rm L} = -t\sum_s\left(\sum_{m_x=-\infty}^{-1}\sum_{m_y=1}^{W_{\rm W}} c^\dagger_{(m_x+1,m_y)s}c_{(m_x,m_y)s} + \sum_{m_x=-\infty}^{0}\sum_{m_y=1}^{W_{\rm W}-1} c^\dagger_{(m_x,m_y+1)s}c_{(m_x,m_y)s} + {\rm H.c.}\right), \tag{3}$$

$$\begin{aligned} H_{\rm R} = -t\sum_s\Bigg(&\sum_{m_x=m_x^R+1}^{\infty}\sum_{m_y=m_y^B}^{m_y^{\rm T}} c^\dagger_{(m_x+1,m_y)s}c_{(m_x,m_y)s} \\ &+ \sum_{m_x=m_x^R+1}^{\infty}\sum_{m_y=m_y^B}^{m_y^{\rm T}-1} c^\dagger_{(m_x,m_y+1)s}c_{(m_x,m_y)s} + {\rm H.c.}\Bigg), \end{aligned} \tag{4}$$

$$\begin{aligned} H_{\rm W} = -\sum_s\Bigg\{&t\sum_{m_xm_y}\left(c^\dagger_{(m_x+1,m_y)s}c_{(m_x,m_y)s} + c^\dagger_{(m_x,m_y+1)s}c_{(m_x,m_y)s}\right) \\ &+ {\rm i}t_{\rm SO}\sum_{s'}\left(c^\dagger_{(m_x+1,m_y)s}(\tau_y)_{ss'}c_{(m_x,m_y)s'} - c^\dagger_{(m_x,m_y+1)s}(\tau_x)_{ss'}c_{(m_x,m_y)s'}\right) + {\rm H.c.}\Bigg\}, \end{aligned} \tag{5}$$

$$H_{\mathrm{T}} = -t_{\mathrm{L}} \sum_{s} \left(\sum_{m_y=1}^{W_{\mathrm{W}}} c^{\dagger}_{(1,m_y)s} c_{(0,m_y)s} + \sum_{m_y=m_y^B}^{m_y^{\mathrm{T}}} c^{\dagger}_{(m_x^R+1,m_y)s} c_{(m_x^R,m_y)s} + \mathrm{H.c.} \right). \tag{6}$$

Here, $c_{(m_x,m_y)s}$ is the electronic annihilation operator at the site $\vec{m} = (m_x, m_y)$ with the spin index $s = \uparrow$ or $\downarrow$. The electronic movement in the transverse direction is restricted by the hard-wall boundary condition adopted in the lateral sides of the wire and the leads. The meanings of m_x^R, m_y^B and m_y^{T} are illustrated in figure 1. With the z-axis set as the quantization direction,

$$\tau_x = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \qquad \tau_y = \begin{pmatrix} 0 & -\mathrm{i} \\ \mathrm{i} & 0 \end{pmatrix} \qquad \text{and} \qquad \tau_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}.$$

The tight-binding Hamiltonian (2) is obtained from the effective mass Hamiltonian (1) by employing the local orbital basis and has been successfully applied to quasi-1D and 2D mesoscopic Rashba structures [9–13, 20]. It is related to the effective mass Hamiltonian (1) via the relations $t = \hbar^2/(2m^*a^2)$, $t_{\mathrm{SO}} = \alpha/(2a)$ and $L_{\mathrm{SO}} = \pi t a/(2t_{\mathrm{SO}})$ with a the lattice spacing. To obtain equation (2), the confining potential V_{conf} is assumed to be zero in the wire and the leads, and infinity outside.

The Green function of the wire can be written as [11, 12, 21, 22]

$$\begin{aligned}\left(G^{-1}(\epsilon)\right)_{\vec{m}\vec{m}',ss'} = {} & \epsilon\delta_{\vec{m},\vec{m}'}\delta_{s,s'} - (H_{\mathrm{W}})_{\vec{m}\vec{m}',ss'} - (\Sigma_{\mathrm{L}}(\epsilon))_{m_y m_y',ss'}\,\delta_{m_x,1}\delta_{m_x',1} \\ & - (\Sigma_{\mathrm{R}}(\epsilon))_{m_y m_y',ss'}\,\delta_{m_x,m_x^R}\delta_{m_x',m_x^R}\end{aligned} \tag{7}$$

where $\vec{m}$ and $\vec{m}'$ can only take sites on the wire. Here, $\Sigma_{\mathrm{L(R)}}$ is the self-energy due to the left (right) semi-infinite ideal leads. The expressions of them can be obtained from the formalism of principal layers in the framework of the surface Green function matching theory with the help of an iterative procedure [21–24]. Once they are known, the coupling function $\Gamma_{\mathrm{L(R)}}$ between the left (right) lead and the wire can be easily obtained as

$$(\Gamma_{\mathrm{L}})_{m_y m_y',ss'} = \mathrm{i}\left(\Sigma_{\mathrm{L}}^r - \Sigma_{\mathrm{L}}^a\right)_{m_y m_y',ss'} \tag{8}$$

with $1 \leqslant m_y, m_y' \leqslant W_{\mathrm{W}}$ and

$$(\Gamma_{\mathrm{R}})_{m_y m_y',ss'} = \mathrm{i}\left(\Sigma_{\mathrm{R}}^r - \Sigma_{\mathrm{R}}^a\right)_{m_y m_y',ss'} \tag{9}$$

with $m_y^B \leqslant m_y, m_y' \leqslant m_y^{\mathrm{T}}$. Here, the variable ϵ is omitted for conciseness of the expressions. Then, the transmission matrix $T^{(\mathrm{LR})}$ can be written as

$$T^{(\mathrm{LR})}_{m_y m_y',ss'} = \sum_{m_y'' m_y''',s''s'''} \left(\sqrt{\Gamma_{\mathrm{L}}}\right)_{m_y m_y'',ss''} \times G^r_{(1,m_y'')(m_x^R,m_y'''),s''s'''} \left(\sqrt{\Gamma_{\mathrm{R}}}\right)_{m_y''' m_y',s'''s'}, \tag{10}$$

where G^r is the retarded Green function of the wire with ϵ replaced by $\epsilon + \mathrm{i}\eta$ in equation (7). To obtain the square root of the matrix $\Gamma_{\mathrm{L(R)}}$, it is first expanded as $\Gamma_{\mathrm{L(R)}} = U^{\dagger}_{\mathrm{L(R)}} D_{\mathrm{L(R)}} U_{\mathrm{L(R)}}$ with $D_{\mathrm{L(R)}}$ a diagonal matrix. Because of the spin degeneracy in the leads, the above four equations can be further simplified.

The transmission coefficient of a spin-s electron incident from the left lead to be transmitted to the right lead as a spin-s' electron is

$$T^{(\mathrm{LR})}_{ss'} = \sum_{m_y=1}^{W_{\mathrm{W}}} \sum_{m_y'=m_y^B}^{m_y^{\mathrm{T}}} \left\| T^{(\mathrm{LR})}_{m_y m_y',ss'} \right\|^2. \tag{11}$$

From the Landauer–Büttiker formula, the conductance of an electron at the Fermi surface ϵ_{F} is

$$\sigma = \sum_{ss'} T^{(\mathrm{LR})}_{ss'}(\epsilon_{\mathrm{F}}) \tag{12}$$

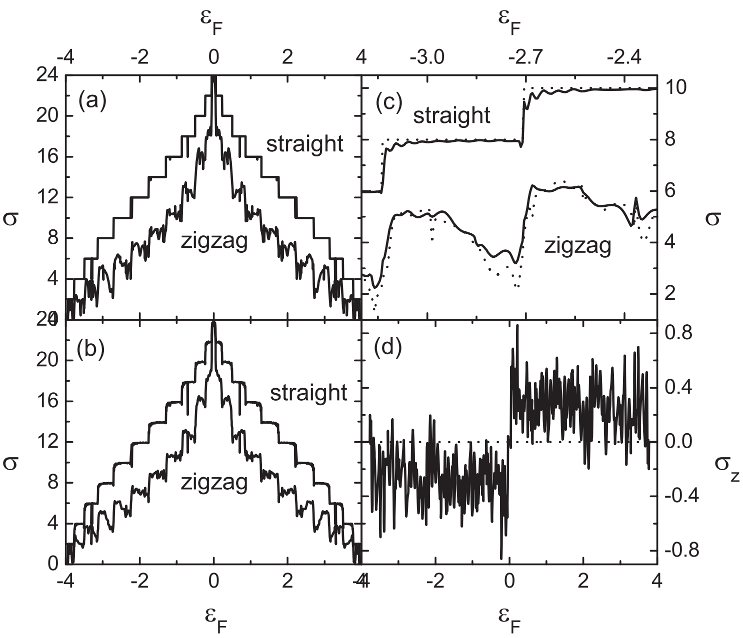

Figure 2. σ–ϵ_F curves for straight and zigzag wires with $t_{SO} = 0$ (a) and 0.1 (b). (c) Details of two conductance steps with $t_{SO} = 0$ (dotted) and 0.1 (solid). (d) σ_z–ϵ_F curves with $t_{SO} = 0.1$ for the straight (dotted) and zigzag (solid) wires. The other parameters are $t_L = t = 1$, $L_W = 60$, $W_W = 12$ and $m_y^T = 36$.

with the unit e^2/h omitted. Here, the temperature is set as zero. Similarly, the spin conductance of the z-component is

$$\sigma_z = \sum_s \left(T_{s\uparrow}^{(LR)}(\epsilon_F) - T_{s\downarrow}^{(LR)}(\epsilon_F) \right), \tag{13}$$

with the unit $\frac{e}{4\pi}$ omitted.

In this method, the self-energy $\Sigma_{L(R)}$ is obtained strictly via numeric calculation, and, consequently, scattering at corners and interfaces between the wire and the leads is treated exactly. With the help of this numerically strict method, the influences of structural symmetry and Rashba SO coupling on σ_z can be studied in detail.

3. Results and discussion

The variations of σ and σ_z with the Fermi energy ϵ_F are plotted in figure 2. In the system of a straight wire, due to confinement in the transverse direction, a series of eigenmodes can be formed in the wire as well as in the leads. Without the Rashba SO coupling, the eigenmodes in the wire are identical to those in the leads. Since t_L is always set as $t_L = t$ in our calculation, no interface scattering occurs in this system, and a series of conductance steps is formed in the σ–ϵ_F curve with each step height being 2. With $t_{SO} \neq 0$, the spin degeneracy of the eigenmodes in the wire is lifted. This leads to interface scattering, but because t_{SO} is set as 0.1 here, its influence is weak (cf figure 2(c)). Despite the lifting of the spin degeneracy, the transparency of an initial spin-up electron is still equal to that of a spin-down electron and $\sigma_z = 0$. With the straight wire turned into a zigzag one, those eigenmodes can still be formed in every straight part of the zigzag structure so that the basic characteristic of conductance steps remains. But unlike the straight wire, σ fluctuates in each step due to corner scattering, and the average heights of those steps are reduced, except that around $\epsilon_F = 0$, which is always $2W_W$. Furthermore, with the Rashba SO coupling introduced into the zigzag one, the transparency of an initial spin-up electron is no longer equal to that of an initial spin-down electron. As a result, the spin conductance of the z-component is nonzero, $\sigma_z \neq 0$ (cf figure 2(d)).

The reason for this great difference in spin transmission can be resolved by analysing the continuum Hamiltonian (1). As we have said, for a straight wire, the Hamiltonian is invariant under reflection with respect to two mirrors. One is along the transverse direction with $(x \rightarrow -x)$, $(p_x \rightarrow -p_x)$ and $(\tau_y, \tau_z \rightarrow -\tau_y, -\tau_z)$, and the other is along the longitudinal direction and has the same transformation relation with x and y exchanged. Here, for simplicity of discussion, the origin of the axis system is moved to the symmetric centre of the structure. According to the former mirror symmetry, $T^{(LR)}_{ss'} = T^{(RL)}_{\bar{s}\bar{s}'}$, whereas from the latter, it can be obtained that $T^{(LR)}_{ss'} = T^{(LR)}_{\bar{s}\bar{s}'}$, which means that $T^{(LR)}_{\uparrow\uparrow} = T^{(LR)}_{\downarrow\downarrow}$ and $T^{(LR)}_{\uparrow\downarrow} = T^{(LR)}_{\downarrow\uparrow}$. This results in zero σ_z. Obviously, to obtain nonzero σ_z, the latter mirror symmetry must be broken. With a straight wire turned into a zigzag one, these two types of mirror symmetry are both destroyed. But in this system, the C_2 rotation symmetry is still reserved; that is, the Hamiltonian (1) is invariant under the transformation $(x, y \rightarrow -x, -y)$, $(p_x, p_y \rightarrow -p_x, -p_y)$ and $(\tau_x, \tau_y \rightarrow -\tau_x, -\tau_y)$, which means that $T^{(LR)}_{ss'} = T^{(RL)}_{ss'}$. Because of the time-reversal symmetry, which corresponds to a transformation $(\vec{r} \rightarrow \vec{r})$, $(\vec{p} \rightarrow -\vec{p})$ and $(\vec{\tau} \rightarrow -\vec{\tau})$, $T^{(LR)}_{ss'} = T^{(RL)}_{\bar{s}'\bar{s}}$. Together with the C_2-rotation symmetry, we obtain the relation $T^{(LR)}_{ss'} = T^{(LR)}_{\bar{s}'\bar{s}}$. This means that $T^{(LR)}_{\uparrow\uparrow} = T^{(LR)}_{\downarrow\downarrow}$, which can be proved by checking the numeric data. However, from the C_2-rotation and time-reversal symmetries, the relation $T^{(LR)}_{\uparrow\downarrow} = T^{(LR)}_{\downarrow\uparrow}$ cannot be derived and $\sigma_z = 0$ is not guaranteed. Of course, this does not mean that nonzero σ_z can be found in any zigzag wire even with spin degeneracy. To generate a spin-polarized longitudinal current in this zigzag Rashba wire, the Rashba SO coupling plays an important role.

Now, we turn our attention to how σ and σ_z vary with t_{SO}. These results are illustrated in figure 3, and they are obtained for wires of width $W_W = 12$ as in figure 2. The influence of the Rashba SO coupling on the particle and spin transmission is weak when $t_{SO} < 0.1$. With t_{SO} further increased, both σ and σ_z oscillate aperiodically. At $\epsilon_F = -3.8$, (cf figures 3(a)–(c)) where only the lowest pair of transverse modes are open, a series of resonant peaks and destructive zeros can be formed in the σ–ϵ_F curves in both straight and zigzag wires, whereas at $\epsilon_F = -3$, (cf figures 3(d)–(f)) where there are four open pairs of transverse modes, neither resonant peak nor destructive zero can be found, which means that resonance or destruction cannot be reached synchronously by transmission through different pairs of modes. A similar dephasing phenomenon has also been found in multichannel mesoscopic Rashba rings [20]. Here, in the straight wire, the spin-dependent quantum interference effect is only caused by interface scattering, whereas in the zigzag one, it is also related to corner scattering. Only in the zigzag wire can nonzero σ_z be found. With respect to σ_z, the most striking difference between the situations with $\epsilon_F = -3.8$ and -3 is that in the former σ_z is seven orders smaller in magnitude than that in the latter, where $\sigma_z \sim 0.1$. Obviously, this difference must be related to the fact that in the former only one pair of modes participate in transmission, whereas in the latter, there are four open pairs.

To clarify this point, we consider the spin transmission through a zigzag Rashba wire of width $W_W = 2$, where there are at most two open pairs of modes. These results are illustrated in figures 4 and 5. As in the wire with $W_W = 12$, a step-like structure can be found in the σ–ϵ_F curve with the step edges located at $\epsilon_F = 1$ and 3, which correspond to the energies where the number of open pairs changes from two to one and from one to zero, respectively. Only in the

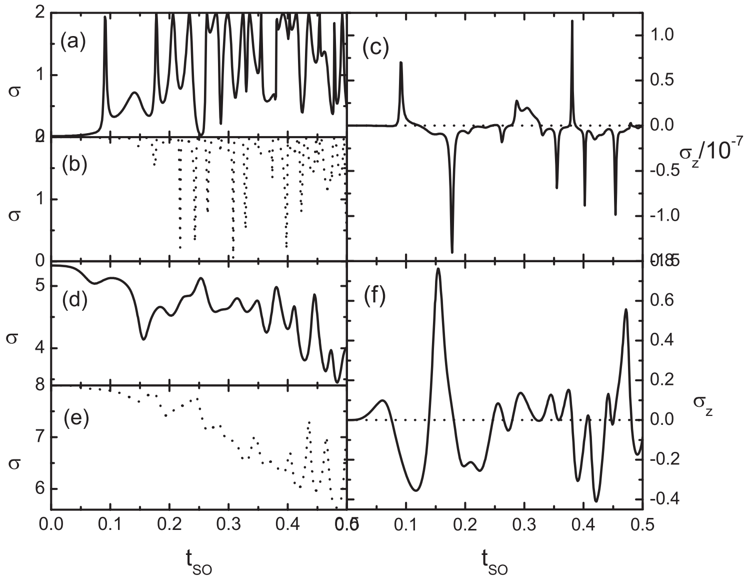

Figure 3. σ–t_{SO} and σ_z–t_{SO} curves at $\epsilon_{\mathrm{F}} = -3.8$ ((a)–(c)) and -3 ((d)–(f)) in zigzag (solid) and straight (dotted) wires. The other parameters are the same as in figure 2.

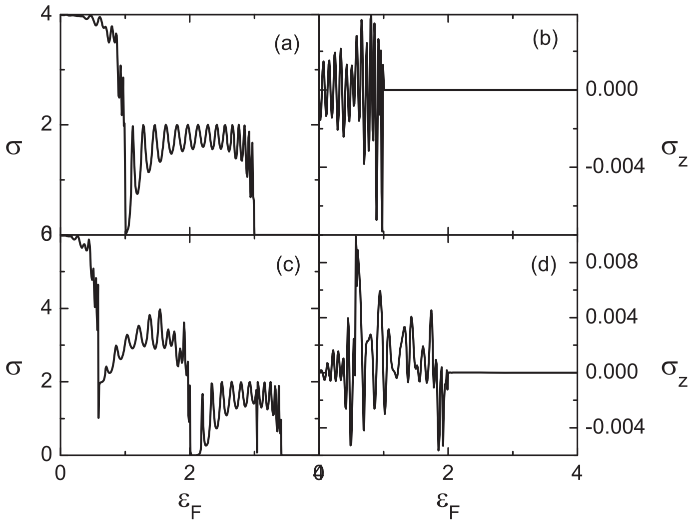

Figure 4. σ–ϵ_{F} and σ_z–ϵ_{F} curves for zigzag wires with $W_{\mathrm{W}} = 2$ ((a) and (b)) and 3 ((c) and (d)) at $t_{\mathrm{SO}} = 0.1$. The other parameters are the same as in figure 2.

region with one open pair of modes can resonant tunnelling be reached. With more than one open pairs, a dephasing phenomenon is unavoidable except at $\epsilon_{\mathrm{F}} = 0$, where the tunnelling through different pairs of modes can reach resonance simultaneously. The spin conductance σ_z oscillates quickly with ϵ_{F} if both pairs are open. With only one open pair, σ_z is greatly reduced, and because it is several orders smaller in magnitude than that in the region with $\epsilon_{\mathrm{F}} < 1$, it looks like zero. (Of course, it is zero only in the region where all pairs of modes are inactive.)

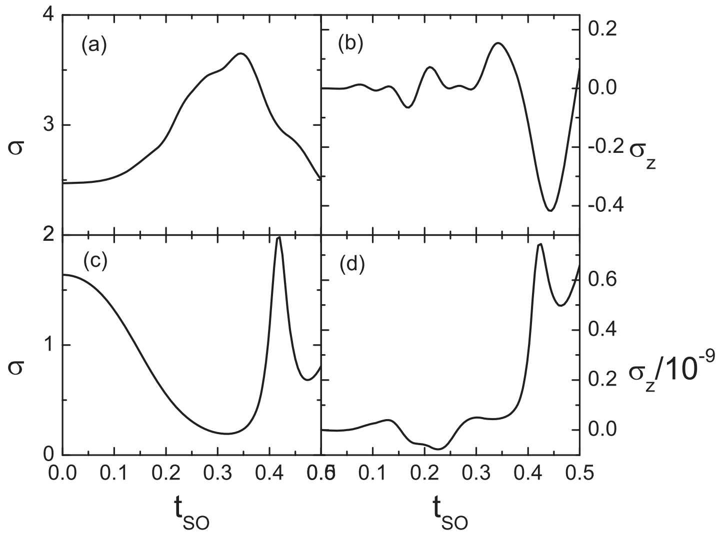

Figure 5. σ–t_{SO} and σ_z–t_{SO} curves for zigzag wires with $W_{\mathrm{W}} = 2$ at $\epsilon_{\mathrm{F}} = 0.9$ ((a) and (b)) and 1.1 ((c) and (d)) with the same other parameters as in figure 2.

Besides the results of the zigzag wire of width $W_{\mathrm{W}} = 2$, those of $W_{\mathrm{W}} = 3$ are also presented in figure 4. They exhibit the same basic characteristics, and one remarkable fact is that with the number of open pairs increased from two to three, a step-like structure cannot be found in the σ_z–ϵ_{F} curve. A similar phenomenon can also be found in the case with $W_{\mathrm{W}} = 12$. Comparing the results of $W_{\mathrm{W}} = 2$ and 3 with those of $W_{\mathrm{W}} = 12$, we can see that the oscillation range of σ_z increases with the wire width, but not proportionally. In figure 4, only the results with $\epsilon_{\mathrm{F}} \geqslant 0$ are illustrated since $\sigma(\epsilon_{\mathrm{F}}) = \sigma(-\epsilon_{\mathrm{F}})$ and $\sigma_z(\epsilon_{\mathrm{F}}) = -\sigma_z(-\epsilon_{\mathrm{F}})$ due to the particle–hole symmetry.

In figure 5, the σ–t_{SO} and σ_z–t_{SO} curves are plotted for a zigzag wire of width $W_{\mathrm{W}} = 2$ at two different Fermi energies: one corresponds to two open pairs of modes and the other to one. In the latter situation, σ can reach resonance whereas it cannot in the former. Meanwhile, in the latter situation, σ_z is generally nine orders smaller in magnitude than that in the former. All of these results mean that mixing between different pairs of modes is crucial for generating experimentally measurable spin-polarized current in these mesoscopic zigzag Rashba wires. Here, the oscillation amplitude of σ_z increases with t_{SO} in both situations, whereas in the wider wire with $W_{\mathrm{W}} = 12$, this monotonic increase is not found. This difference can be seen more clearly by comparison of figures 3(c) and 5(d), where only one pair of modes is open. This demonstrates the influence of those unoccupied modes on transport properties.

In this paper, a mesoscopic zigzag Rashba structure is presented to show a spin-polarized current can be generated along the longitudinal direction in Rashba systems. From symmetry analysis, to have nonzero σ_z, it is necessary only to break the mirror symmetry with respect to the longitudinal direction, and many other structures, even those with left–right symmetry, satisfy this condition. For example, in a mesoscopic Rashba wire with a triangular structure, although the left–right symmetry is still reserved, the mirror symmetry with respect to the longitudinal direction is broken. From the left–right and time-reversal symmetries, the relation $T^{(\mathrm{LR})}_{ss'} = T^{(\mathrm{LR})}_{s's}$ is obtained so that $T^{(\mathrm{LR})}_{\uparrow\downarrow} = T^{(\mathrm{LR})}_{\downarrow\uparrow}$, but since the relation $T^{(\mathrm{LR})}_{\uparrow\uparrow} = T^{(\mathrm{LR})}_{\downarrow\downarrow}$ cannot be guaranteed in this structure, it is still possible to find nonzero σ_z if Rashba SO coupling is introduced.

In all of the above investigation, the influence of impurity scattering is not taken into account. With one impurity randomly located on the wire, any structural symmetry will be destroyed no matter whether the wire is a straight or zigzag one, and the corresponding σ_z is generally nonzero. But usually, the wire contains a lot of impurities, and they are distributed randomly. In this random system, what we are interested in are ensemble-averaged variables. In the process of ensemble averaging, the structural symmetry is recovered. As a result, the ensemble-averaged $\langle\sigma_z\rangle$ is zero for a straight wire, whereas it is nonzero for a zigzag one, as discussed above, although its specific value is depressed by the randomness.

4. Summary

In summary, ballistic transport through a mesoscopic zigzag Rashba wire is studies via the Green function technique, and the transmission of particles and spins is investigated with an unpolarized charge current injected. These results are compared with those of a straight wire. The symmetry analysis shows that in Rashba systems, only structures without mirror symmetry with respect to the longitudinal direction are possible for generating nonzero σ_z. This necessary condition is satisfied in the zigzag wire. In this system, the Rashba SO coupling lifts the spin degeneracy, and the scattering at corners, as well as at interfaces, mixes the transverse modes. The ballistic transport properties are influenced by the spin-dependent quantum interference effect. Mixing between different pairs of modes is crucial for spin transmission, and only if more than one pair of modes participate in transport can experimentally measurable σ_z be found. This nonzero σ_z varies with t_{SO}. Our theoretical results provide one way to generate an artificially controllable spin-polarized current in mesoscopic Rashba systems.

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