Electronic transport through a quantum ring coupled to ferromagnetic leads*

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We study the spin-dependent transport through a one-dimensional quantum ring with taking both the Rashba spin–orbit coupling (RSOC) and ferromagnetic leads into consideration. The linear conductance is obtained by the Green's function method. We find that due to the quantum interference effect arising from the RSOC-induced spin precession phase and the difference in travelling phase between the two arms of the ring, the conductance becomes spin-polarized even in the antiparallel magnetic configuration of the two leads, which is different from the case in single conduction channel system. The linear conductance, the spin polarization and the tunnel magnetoresistance are periodic functions of the two phases, and can be efficiently tuned by the structure parameters.

Keywords: quantum ring, rashba spin–orbit interaction, ferromagnetic leads

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1. Introduction

One-dimensional quantum ring is a subject frequently studied in solid-state systems, due to its interesting quantum interference effect.[1−3] Especially in recent years, various ring-type models were extensively investigated to control the electron spin degree of freedom, most of which took the Rashba spin–orbit coupling (RSOC) into account.[4−15] The RSOC arises from the lack of structural inversion symmetry in growth direction, and is essentially the influence of an electric field on a moving spin, thereby providing an electrical spin control method.[12−14,16−18] Many recent experimental and theoretical studies indicate that the spin polarization based on the RSOC can reach as high as 100%[15,19] or infinite[20−30] by tuning external electrical field or gate voltages. For example, spin-up and spin-down electrons were proposed to be split individually to different drain leads from a spin-unpolarized source lead in a Rashba ring.[19,27] A pure spin current, which means that equal numbers of spin-up and spin-down electrons are moving in the opposite directions, can also be generated in similar devices.[28−30]

Although the spin control efficiency in the above-mentioned studies is quite high and attractive, the real materials in nowadays commercial spintronics devices are mainly made of magnetic ones, relying on the giant magnetoresistance or tunnel magnetoresistance (TMR) effects.[31−33] In view of these, to study the combination of quantum interference effect and the ferromagnetic electron reservoirs should be of importance, which is the motivation of the present paper. In fact, such a topic has been addressed very recently. The current–voltage and the TMR characteristics have been investigated in a quantum ring with two quantum dots (QDs) inserted in its arms.[34−36] The TMR is defined as $(G_{\rm P} - G_{\rm AP})/G_{\rm AP}$, where $G_{\rm P(AP)}$ is the total conductance when the magnetic configuration of the device is in parallel (antiparallel) alignment, and is a key physical quantity in the designing of spintronics device. Previous studies showed that TMR was sensitive to the quantum interference effect, with both its signs and magnitude being adjustable.[34−36]

In the present paper, we study the properties of the spin-dependent linear conductance and the TMR in a one-dimensional ballistic quantum ring coupled to two ferromagnetic leads (see Fig. 1). It should be noted that the device under consideration here is much simpler than those in Refs. [34]–[36], since there are no QDs in our ring arms. Actually, in most of the above-mentioned studies, the QDs were considered, since they were of important structures in low-dimensional solid-state systems,[37−43] but they are neglected here for the sake of device simplicity. If

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there is RSOC in the ring arms, spin-up and spin-down electrons will acquire different phase factors, the so-call spin precession phases, when tunneling through the system.[12−14,16] This phase, combined with the travelling phase difference due to unequal lengths of arms l_1 and l_2,[44] will have a significant influence on the spin-related transport.

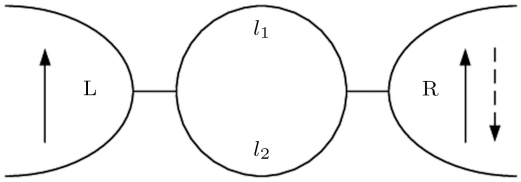


Fig. 1. System of a quantum ring coupled to the left and right ferromagnetic leads. The lengths of the upper and the lower arms are respectively l_1 and l_2.

2. Model and method

The system with a quantum ring coupled to the left and the right ferromagnetic leads can be described by the following Hamiltonian:[44,45]

$$H = \sum_{k,\sigma,\beta=\mathrm{L,R}} \varepsilon_{k\beta\sigma} c^{\dagger}_{k\beta\sigma} c_{k\beta\sigma} + \sum_{\sigma,n=1,2} \varepsilon_0 d^{\dagger}_{n\sigma} d_{n\sigma} + \sum_{\beta,\sigma,n} \Big(V_{\beta n\sigma} c^{\dagger}_{k\beta\sigma} d_{n\sigma} + \mathrm{H.c.} \Big), \quad (1)$$

where $c^{\dagger}_{k\beta\sigma}$ ($c_{k\beta\sigma}$) is the creation (annihilation) operator of the electrons with momentum k, spin index σ ($\sigma = \uparrow, \downarrow$ or $\sigma = \pm 1$) and energy $\varepsilon_{k\beta\sigma}$ in lead β; $d^{\dagger}_{n\sigma}$ ($d_{n\sigma}$) creates (annihilates) an electron in the upper ($n = 1$) and the lower ($n = 2$) arms of the ring with the same on-site energy ε_0 for a uniform ring; $V_{\beta n\sigma}$ describes the dot-lead tunneling coupling (its energy-dependence is neglected for the sake of simplicity); $V_{\mathrm{L}1\sigma} = V_{\mathrm{L}1} \mathrm{e}^{\mathrm{i}\,\theta}$, $V_{\mathrm{L}2\sigma} = V_{\mathrm{L}2}$, and $V_{\mathrm{R}n\sigma} = V_{\mathrm{R}n} \mathrm{e}^{-\mathrm{i}\,\sigma\phi_n}$, with θ being the travelling phase difference for unequal L_1 and L_2,[42] which acts as an external magnetic flux, $\phi_n = \alpha_n m L_n/\hbar^2$ being the RSOC-induced spin precession phase factor in the nth ring arm,[9] with α_n and m being the gate voltage tunable RSOC strength[18] and the electron effective mass, respectively.

The spin-dependent linear conductance at zero temperature is obtained from the Landauer–Büttiker formula as[27]

$$\mathrm{G}_{\sigma}(\varepsilon) = \frac{e^2}{h} \mathrm{Tr}\{ \mathbf{G}^{a}_{\sigma}(\varepsilon) \mathbf{\Gamma}^{\mathrm{R}}_{\sigma} \mathbf{G}^{r}_{\sigma}(\varepsilon) \mathbf{\Gamma}^{\mathrm{L}}_{\sigma} \}, \quad (2)$$

where the 2×2 matrix $\mathbf{G}^{r(a)}_{\sigma}(\varepsilon)$ is the retarded (advanced) Green's function of the ring. The matrix $\mathbf{\Gamma}^{\beta}_{\sigma}$ ($\beta = \mathrm{L,R}$) is the tunnel-coupling strength between the ring and the leads, and is given by

$$\mathbf{\Gamma}^{\mathrm{L}}_{\sigma} = \begin{bmatrix} \Gamma^{\mathrm{L}}_{11\sigma} & \sqrt{\Gamma^{\mathrm{L}}_{11\sigma}\Gamma^{\mathrm{L}}_{22\sigma}}\mathrm{e}^{\mathrm{i}\,\theta} \\ \sqrt{\Gamma^{\mathrm{L}}_{22\sigma}\Gamma^{\mathrm{L}}_{11\sigma}}\mathrm{e}^{-\mathrm{i}\,\theta} & \Gamma^{\mathrm{L}}_{22\sigma} \end{bmatrix}, \quad (3)$$

$$\mathbf{\Gamma}^{\mathrm{R}}_{\sigma} = \begin{bmatrix} \Gamma^{\mathrm{R}}_{11\sigma} & \sqrt{\Gamma^{\mathrm{R}}_{11\sigma}\Gamma^{\mathrm{R}}_{22\sigma}}\mathrm{e}^{-\mathrm{i}\,\sigma\phi} \\ \sqrt{\Gamma^{\mathrm{R}}_{22\sigma}\Gamma^{\mathrm{R}}_{11\sigma}}\mathrm{e}^{\mathrm{i}\,\sigma\phi} & \Gamma^{\mathrm{R}}_{22\sigma} \end{bmatrix}, \quad (4)$$

where $\phi = \phi_1 - \phi_2$. The matrix elements in the above two equations are defined as $\Gamma^{\beta}_{mn\sigma} = 2\pi V_{\beta m\sigma} V^{*}_{\beta n\sigma} \rho^{\beta}_{\sigma}$. ρ^{β}_{σ} is the density of states in lead β (its energy-dependence is neglected), which describes the ferromagnetism of the leads. For simplicity we assume symmetric couplings between the ring and the leads with $V_{\mathrm{L}1} = V_{\mathrm{L}2} = V_{\mathrm{R}1} = V_{\mathrm{R}2}$, so we have $\Gamma^{\beta}_{11\sigma} = \Gamma^{\beta}_{22\sigma}$. Then the ferromagnetism of the leads is expressed by introducing a spin-polarization parameter p_{β} for the two leads,[34−36,46] with $p_{\beta} = (\Gamma^{\beta}_{nn\uparrow} - \Gamma^{\beta}_{nn\downarrow})/(\Gamma^{\beta}_{nn\uparrow} + \Gamma^{\beta}_{nn\downarrow})$, in which $n = 1, 2$. In this paper, we consider the parallel ($p_{\mathrm{L}} = p_{\mathrm{R}} = p$) and antiparallel ($p_{\mathrm{L}} = -p_{\mathrm{R}} = p$) configurations of the two leads. In both cases we have $\Gamma^{\mathrm{L}}_{nn\uparrow(\downarrow)} = (1 \pm p)\Gamma$, while $\Gamma^{\mathrm{R}}_{nn\uparrow(\downarrow)} = (1 \pm p)\Gamma$ is for the parallel case and $\Gamma^{\mathrm{R}}_{nn\uparrow(\downarrow)} = (1 \mp p)\Gamma$ is for the antiparallel case, with $\Gamma = (\Gamma^{\beta}_{nn\uparrow} + \Gamma^{\beta}_{nn\downarrow})/2$.

The Green's functions in Eq. (2) is obtained from the Dyson equation as

$$\mathbf{G}^{r(a)}_{\sigma}(\varepsilon) = \mathbf{g}^{r(a)}_{\sigma}(\varepsilon) + \mathbf{g}^{r(a)}_{\sigma}(\varepsilon) \mathbf{\Sigma}^{r(a)}_{\sigma} \mathbf{G}^{r(a)}_{\sigma}(\varepsilon), \quad (5)$$

where $\mathbf{g}^{r(a)}_{\sigma}(\varepsilon)$ is the retarded (advanced) Green's function for the uncoupled ring, with its matrix elements given by

$$\mathrm{g}^{r}_{mn\sigma}(\varepsilon) = \mathrm{g}^{a*}_{mn\sigma}(\varepsilon) = \frac{1}{\varepsilon - \varepsilon_0 + \mathrm{i}\,0^{+}} \delta_{mn}. \quad (6)$$

The self-energy in Eq. (5) is $\mathbf{\Sigma}^{r(a)}_{\sigma} = \mp \mathrm{i}(\mathbf{\Gamma}^{\mathrm{L}}_{\sigma} + \mathbf{\Gamma}^{\mathrm{R}}_{\sigma})/2$.

After solving the Green's functions, the spin-dependent linear conductances for parallel ($\mathrm{G}_{\mathrm{P}\sigma}$) and antiparallel ($\mathrm{G}_{\mathrm{AP}\sigma}$) configurations can be written explicitly as

$$G_{P\sigma}(\varepsilon)=\frac{e^2}{h}\frac{4\Gamma^2(1+\sigma p)^2(\varepsilon-\varepsilon_0)^2\cos^2\dfrac{\theta+\sigma\phi}{2}}{[(\varepsilon-\varepsilon_0)^2-\Gamma^2(1+\sigma p)^2\sin^2\dfrac{\theta+\sigma\phi}{2}]^2+4\Gamma^2[(1+\sigma p)(\varepsilon-\varepsilon_0)]^2} \tag{7}$$

and

$$G_{AP\sigma}(\varepsilon)=\frac{e^2}{h}\frac{4\Gamma^2(1-p^2)(\varepsilon-\varepsilon_0)^2\cos^2\dfrac{\theta+\sigma\phi}{2}}{[(\varepsilon-\varepsilon_0)^2-\Gamma^2(1-p^2)\sin^2\dfrac{\theta+\sigma\phi}{2}]^2+4\Gamma^2(\varepsilon-\varepsilon_0)^2}. \tag{8}$$

Note that if $p=0$ (the normal metal leads case), the conductances in the above two equations are just the same as those in Ref. [15] of a ring with two noninteracting QDs. This is because when the QDs are free from Coulomb interaction, they just provide a discrete on-site energy, which has no essential difference from the result of the present model, even though the experimental requirement of the present device is much lower than that of Refs. [15], [34–36], [43].

Before presenting numerical results, we first discuss some general characteristics of the linear conductances given in Eqs. (7) and (8). (i) In the absence of the RSOC, the antiparallel configuration conductance $G_{AP\sigma}$ is spin-independent while $G_{P\sigma}$ depends on spin, which is a usual case.[34–36] Even in the presence of RSOC but with zero θ, $G_{AP\sigma}$ can never become spin-polarized because of the phase locking effect.[12–15] But in the case where both θ and ϕ are nonzeros, $G_{AP\sigma}$ may depend on electron spin. (ii) Due to the existence of $\cos^2\frac{\theta+\sigma\phi}{2}$ in the numerators of Eqs. (7) and (8), both $G_{P\sigma}$ and $G_{AP\sigma}$ are periodic functions of phases θ and ϕ. (iii) Both $G_{P\sigma}$ and $G_{AP\sigma}$ have two Breit–Wigner resonances centred at $\varepsilon_0\pm\Gamma(1+\sigma p)\sin\frac{\theta+\sigma\phi}{2}$ and $\varepsilon_0\pm\Gamma\sqrt{(1-p^2)}\sin\frac{\theta+\sigma\phi}{2}$, respectively. It indicates that the resonance positions of the conductances are affected by not only the phase factors but also the ferromagnetism of the leads. (iv) The resonance heights of the conductances for the parallel and antiparallel configurations are different from each other, and they are $G_{P\sigma}=\frac{e^2}{h}\cos^2\frac{\theta+\sigma\phi}{2}$ and $G_{AP\sigma}=\frac{e^2}{h}(1-p^2)\cos^2\frac{\theta+\sigma\phi}{2}$, respectively.

3. Numerical results

In this section, we give numerical results about the linear conductance, spin-polarization $\eta_{P(AP)}=(G_{P(AP)\uparrow}-G_{P(AP)\downarrow})/G_{P(AP)}$ and the TMR. Throughout the paper we choose $\Gamma=1$ as the energy unit and set the constants $e=\hbar=k_B=1$ and the spin polarization of the leads $p=0.6$. Since the on-site energy ε_0 just shifts the resonance position, we also set it to be zero throughout the paper. The linear conductances and the TMRs each as a function of the Fermi level ε_F are shown in Fig. 2 for different values of θ and fixed $\phi=0$. As indicated previously, the system conductance for the antiparallel magnetic alignment $G_{AP\sigma}$ is spin-independent in the absence of RSOC, and then only $G_{AP}=G_{AP\uparrow}+G_{AP\downarrow}$ is plotted there. For $\theta=0$ as shown in Fig. 2(a), both $G_{P\sigma}$ and G_{AP} have a resonance centred at ε_0, where $G_{P\uparrow}$ and $G_{P\downarrow}$ reach their quantum limit e^2/h. While G_{AP} at the resonance energy position is $2(1-p^2)e^2/h$ (thick solid line), which is smaller than $G_P=2e^2/h$ (thin solid line). Apart from the resonance energy position, $G_{P\uparrow}$ is larger than $G_{P\downarrow}$ since they are approximately proportional to $(1+p)^2$ and $(1-p)^2$, respectively. When $\theta\neq0$ as shown by Figs. 2(b) and 2(c), the linear conductances are significantly influenced by the quantum interference effect, and have two distinct characteristics as compared with those in Fig. 2(a). First, the Breit–Wigner resonance at ε_0 is now split into two, with positions centred at $\varepsilon_0\pm\Gamma(1+\sigma p)\sin\frac{\theta+\sigma\phi}{2}$ for parallel magnetic alignment and $\varepsilon_0\pm\Gamma\sqrt{(1-p^2)}\sin\frac{\theta+\sigma\phi}{2}$ for the antiparallel case. For $\varepsilon_F=\varepsilon_0$, the conductances are reduced to zero due to the level attraction effect.[15,47] Second, the magnitudes of the conductances are reduced by the modulation of the factor $\cos^2(\theta/2)$. Figure 2(d) shows the properties of the TMR. For $\theta=0$, the curve has a broad resonance centred at ε_0 and two dips, which are induced by the different resonance widths of $G_{P\uparrow}$ and $G_{P\downarrow}$ [see the dash and the dot lines in Fig. 2(a)]. For $\theta\neq0$, the magnitude and the resonance or dip positions of the TMR are significantly changed due to the characteristics of the linear conductances. Since the on-site energies of the two arms of the ring are the same, the linear conductance and the corresponding TMR are both symmetrical functions of Fermi energy with respective to ε_0. Moreover, the TMR value in

Fig. 2(d) is always positive because of $G_P > G_{AP}$ given in Figs. 2(a)–2(c).

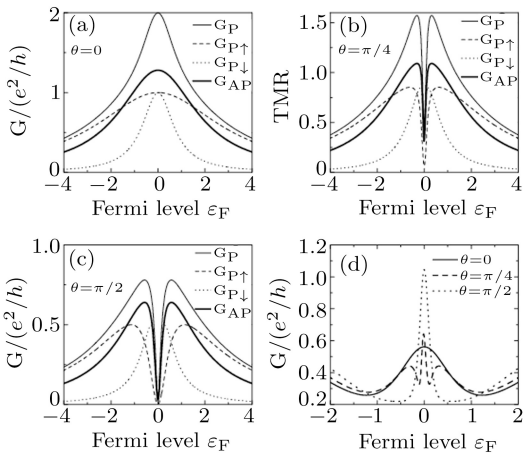


Fig. 2. Linear conductances [(a)–(c)] and the TMRs (d) each as a function of the Fermi level ε_F for zero ϕ and different values of θ.

We now study in Fig. 3 the linear conductance and the TMR each varying as a function of θ for different values of Fermi energy ε_F. It is found that $G_{P(AP)\sigma}$ and the TMR are both 2π-periodic functions of θ. When the Fermi energy ε_F approaches to the on-site energy of the ring ε_0 (Fig. 3(a)), the conductances of both parallel and antiparallel magnetic configurations reach their maximum values for $\theta = 2n\pi$, where $n = 0, 1, 2, \cdots\cdots$. While they are greatly suppressed for other values of θ, which can also be seen from Figs. 2(a) to 2(c). This indicates that the present device may operate as a conductance switcher based on θ. Experimentally, the value of θ, which comes from the asymmetry between the two ring arms, can be adjusted by modulating the electron wave-vectors in one of the ring arms by a gate voltage.[44] When the difference between ε_F and ε_0 increases, the conductance blockade region in θ reduces as shown in Figs. 3(b) and 3(c). For $\varepsilon_F = 0.01$, the TMR is a constant in the region of θ where the conductance is suppressed [see the solid line in Fig. 3(d)]. With the increase of ε_F, this platform is narrowed and lowered (the dash line). For $\varepsilon_F = 0.5$ as indicated by the dot line, the TMR value is even suppressed as compared with the cases of $\varepsilon_F = 0.01$ and $\varepsilon_F = 0.1$.

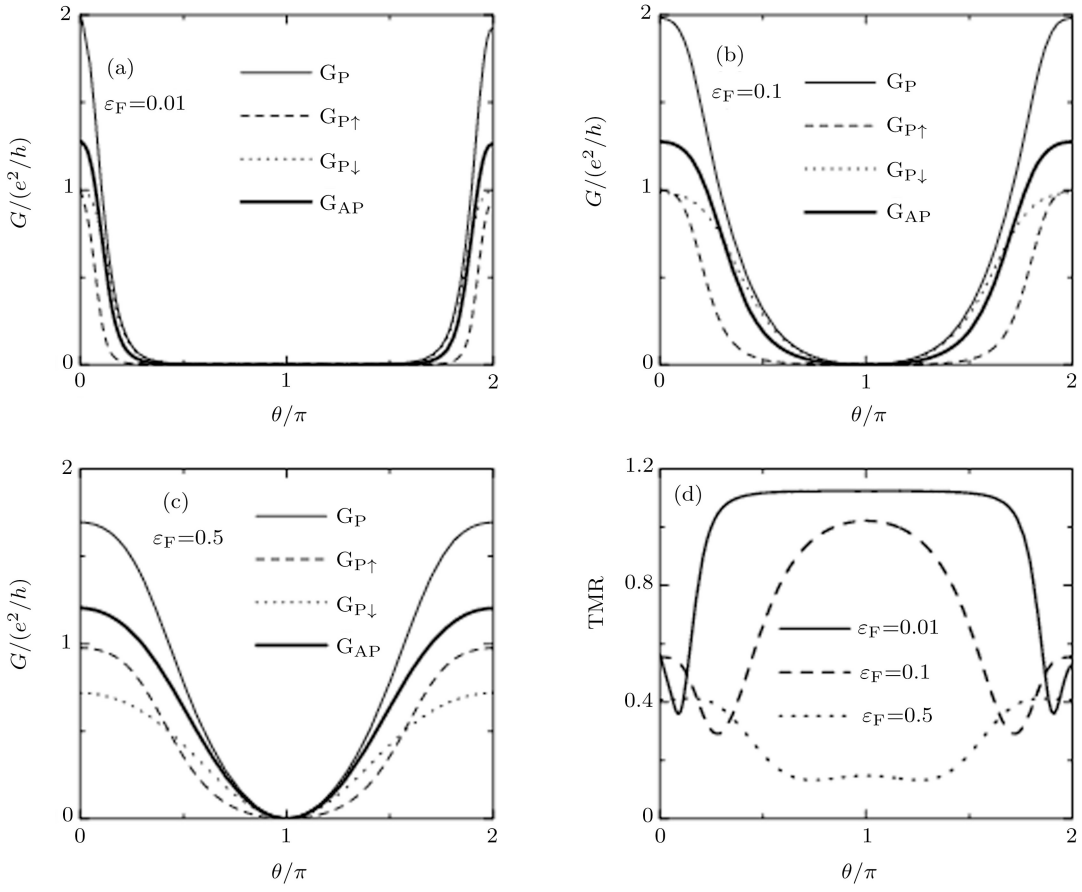


Fig. 3. Linear conductances ((a)–(c)) and the TMRs (d) each as a function of θ for different values of ε_F and $\phi = 0$.

From Eqs. (7) and (8) one can see that the RSOC-induced spin precession phase ϕ can play the same role as θ for totally symmetrical ring arms ($\theta = 0$). For an uniform ring, ϕ can reach as large as about 2π if the ring lengths $l_1 = l_2 \approx 0.1$ μm and the RSOC strengths $\alpha_1 = \alpha_2 \approx 1.05 \times 10^{-11}$ eV·m, which is tunable by gate voltage in experiment.[18] Since the behaviours of the linear conductance as functions of Fermi level and θ are clear, we now focus our attention on the spin polarization and the TMR as functions of θ for different values of ϕ. It can be seen from Fig. 4(a) that the conductance for antiparallel configuration becomes spin-polarized when $\phi = \pi/2$ and $3\pi/2$, while it is independent of spin when $\phi = \pi$. This can be deduced from Eq. (8) that for $\phi = \pi$, since $\cos^2 \frac{\theta+\pi}{2} = \cos^2 \frac{\theta-\pi}{2} = \sin^2 \frac{\theta}{2}$, we have $G_{AP\uparrow} = G_{AP\downarrow}$. The spin polarizations of $\phi = \pi/2$ (the solid line) and $3\pi/2$ (the dot line) are reciprocal to each other, which originates from the different interference effects of the two spin components of electrons. As shown in Fig. 4(b), the spin polarization for the parallel configuration η_P is quite different from that for the antiparallel case η_{AP}. First, η_P values for $\phi = \pi/2$ and $3\pi/2$ are not always opposite to each other any more. Second, the spin polarization η_P for $\phi = \pi$ is finite as compared with the antiparallel case η_{AP}, because of the ferromagnetism of the leads. It is shown that both η_P and η_{AP} can be tuned in a wide range, for example, from -1 to 1, by varying either θ or ϕ. This characteristic may serve as a spin filter that allows the transportation of only one direction of electron spin. We find in Fig. 4(d) that the TMR now can be either positive or negative in the case of nonzero ϕ, which is different from those in Figs. 2 and 3.

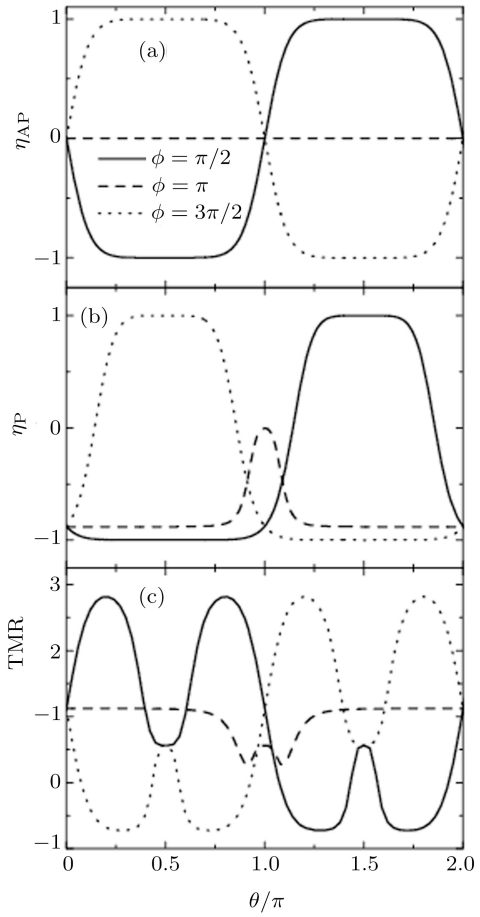


Fig. 4. Spin polarizations ((a) and (b)) and the TMRs (c) each as a function of θ for different values of ϕ and $\varepsilon_F = 0.01$.

4. Conclusion

We have investigated the characteristics of the linear conductance, the spin polarization and the TMR in a one-dimensional quantum ring coupled to ferromagnetic leads. Both the difference in travelling phase between the two ring arms and the Rashba spin precession phase are taken into consideration. It is found that these physical quantities are all 2π-periodic functions of the above two phases. The linear conductances in not only parallel but also antiparallel magnetic configurations can become spin-polarized due to the co-action of the phase factors. Both the signs and the magnitudes of the spin polarization and the TMR can be adjusted with the help of either travelling phase difference or the spin precession phase, both of which can be tuned by varying gate voltage. Under certain conditions, this device can serve as a conductance switcher or a spin filter.

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