Spin-polarized transport induced by spin-pumping in a Rashba ring

2009 J. Phys.: Condens. Matter 21 485304

(http://iopscience.iop.org/0953-8984/21/48/485304)

IOP Publishing Journal of Physics: Condensed Matter
J. Phys.: Condens. Matter **21** (2009) 485304 (7pp) doi:10.1088/0953-8984/21/48/485304

Spin-polarized transport induced by spin-pumping in a Rashba ring

F Liang[1], Y H Yang[1], J Wang[1,2] and K S Chan[2]

[1] Department of Physics, Southeast University, Nanjing 210096, People's Republic of China
[2] Department of Physics and Materials Science, City University of Hong Kong, Tat Chee Avenue, Kowloon, Hong Kong, People's Republic of China

E-mail: lfgy812512@gmail.com and apkschan@cityu.edu.hk

Received 3 September 2009, in final form 14 October 2009
Published 6 November 2009
Online at stacks.iop.org/JPhysCM/21/485304

Abstract
The Keldysh Green's function method is employed to study spin-dependent electron transport through a Rashba ring with a quantum dot (QD) embedded in one of its arms. Zero charge bias is applied on the system while a rotating magnetic field is considered in the QD to pump pure spin current. The Rashba spin–orbital coupling (RSOC) can cause a spin precession phase of the electron passing through the ring, so that the quantum interference in the ring can lead to a spin-polarized charge current flowing in the leads and the arm without a QD, whereas only pure spin current is flowing in the other arm with a QD. It is shown that for low frequency ω of the rotating magnetic field, the pumped charge current is proportional to ω unlike the charge current produced by mono-parametric quantum charge pumping, which is usually proportional to ω^2. Moreover, the magnitude, the direction, as well as the spin-polarization of the charge current can be controlled by tuning the device parameters such as the QD energy level, the RSOC strength, and the strength of the electron tunneling between the leads and the QD. Hence the studied device may serve as a generating source for tunable spin-polarized current in the spintronics field.

1. Introduction

Since Datta and Das [1] theoretically proposed the possibility of the spin transistor almost two decades ago, spin-related transport phenomena in semiconductor nanostructures has attracted more and more attention [2–5]. This has resulted in the emergence of semiconductor spintronics that aims to develop spin-based electronic devices utilizing not only the charge but also the spin degrees of freedom of an electron. Because the spin-based electronic devices have many advantages over the traditional ones [6] such as longer coherent lifetime, faster data proceeding speed, and lower electric power consumption, the subject of generating a spin-polarized current through a mesoscopic system has become one of the key issues in the field of spintronics and has prompted intense interest in recent years.

At present, many methods for producing spin current have been proposed in the literature [7–18]. For instance, the spin-Hall effect exploiting the intrinsic spin–orbit coupling in nonmagnetic semiconductors [7, 8] offers an effective way to generate a dissipationless spin current. The generic magnetic means for producing spin current mainly includes spin injection from magnetic materials [9–12] and spin pumps [13–18] driven by time-dependent external fields. Wang *et al* proposed a spin field effect transistor (SFET) [13] which operates with a rotating magnetic field (RMF). While the original spin-degenerate electron levels in the scattering region are split by the static longitudinal component of the RMF, the rotating transverse component of the RMF couples the two Zeeman levels and makes transition between them possible. An electron can then tunnel into the spin-down level from an external lead and then it will tunnel out of the scattering region with up-spin by a spin-flip process due to the transverse RMF, leaving behind a spin-down hole, which will subsequently be filled with an incoming electron from the external lead again. This cyclic process leads to a spin current flowing in the device. As the source–drain bias is zero, only a pure spin current is pumped out by RMF without any accompanying charge current.

Recently, much attention has also been paid to the spin-dependent transport phenomena in various kinds of ring-type or two-path semiconductor structures, such as an Aharonov–Bohm (AB) ring [19–21], a Stern–Gerlach ring [22], an Aharonov–Bohm–Casher ring coupled to a QD [23], a ring with serially coupled double quantum dots [24], an AB ring with a QD embedded in one of its arms [25], a two-path semiconductor device [26], and so on. In these systems, the presence of the Rashba spin–orbit coupling (RSOC) can cause the electrons moving through the ring to acquire a spin-dependent phase (Aharonov–Casher (AC) phase). The resultant spin-dependent phase difference alone or together with some other nondispersive phase (e.g. an AB phase induced by an external magnetic field perpendicular to the ring) can result in an interesting spin-interference effect that makes the system conductance spin-polarized, i.e. a spin-polarized current can be achieved with an unpolarized injecting charge current. By using a gate voltage to vary the RSOC strength [27, 28], both the magnitude and spin-polarization of the generated current can be tuned conveniently. While many works have been dedicated to exploiting the spin-pump or spin-interference effect to realize the generation of a spin-polarized current, the combined effect of the two effects on the spin-polarized transport in a nanostructure is still lacking attention. In the present work, we demonstrate that the combined effect of spin-pump and spin-interference can give rise to a controllable spin-polarized charge current in a one-dot Rashba ring device at a zero bias. In the proposed device, a rotating magnetic field is applied to the quantum dot (QD) embedded in one of the arms of the ring, which plays the role of a spin-pump source. With the help of the Keldysh Green's function technique, we show that a spin-polarized charge current can flow in the system without any external bias, because the quantum spin-interference effect in the ring breaks the symmetry of up-spin and down-spin current of the pure spin current generated by the RMF. The spin-polarized charge current is proportional to the driving frequency ω and its magnitude, direction, and spin-polarization can be modulated by tuning the device parameters such as the energy level in the QD, the RSOC constant, as well as the coupling strength between the leads and ring. Of particular interest is that a fully spin-polarized charge current (one spin species current is suppressed) can be achieved, i.e. spin injection with 100% efficiency can be realized in our proposed device.

This paper is organized as follows. In section 2, we introduce the model and present the formulae to calculate the generated spin-polarized current in the presence of the RMF. In section 3, the corresponding numerical results are given. In section 4, a brief conclusion is presented.

2. Model and formulation

We consider a two-terminal device in which a Rashba ring in the x–y plane is connected with the left and right normal leads, as schematically shown in figure 1, a noninteracting QD sits on one arm of the ring (say upper arm), and an RMF is applied on it. The RMF rotates with angular frequency ω around the z-axis with a tilt angle θ and is expressed as: $\mathbf{B} = (B_0 \sin\theta \cos\omega t, B_0 \sin\theta \cos\omega t, B_0 \cos\theta)$ with B_0 being the constant magnetic field strength. The z component $B_z = B_0 \cos\theta$ gives a Zeeman split to the original spin-degenerate dot levels. The transverse component $(B_0 \sin\theta \cos\omega t, B_0 \sin\theta \cos\omega t)$ in the x–y plane provides a spin-flip mechanism between the two Zeeman levels. A

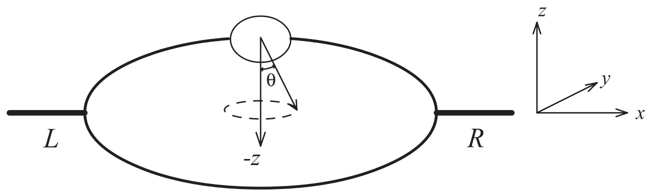

Figure 1. The scheme of the device considered in this work, a one-dot Rashba ring coupled to two leads. The ring, which lies in the x–y plane, contains a dot embedded in its upper arm and the RSOC is considered in the ring. An RMF, which here plays the role of a spin-pump source, is applied to the dot.

uniform RSOC is considered in the ring and it can induce a spin-dependent phase shift $\sigma\varphi$ (AC phase) of electron tunneling through the ring with $\varphi = -\alpha m L/\hbar^2$ with m the effective mass of an electron, α the RSOC strength, and L the length of the ring. In our model, the phase shift $\sigma\varphi$ from RSOC is considered between the QD and right lead without loss of generalization, which can certainly denote a hybrid ring with nonuniform RSOC. The one-dot Rashba ring device with the RMF can be described by the following Hamiltonian:

$$H = \sum_{\alpha=\mathrm{L,R}} H_\alpha + H_\mathrm{D} + H'(t) + H_\mathrm{T}, \tag{1}$$

with

$$H_\alpha = \sum_{k\sigma} \epsilon_{\alpha k} C^\dagger_{\alpha k\sigma} C_{\alpha k\sigma}, \tag{2}$$

$$H_\mathrm{D} = \sum_\sigma (\varepsilon_d + \sigma\mu_\mathrm{B} B_z) d^\dagger_\sigma d_\sigma, \tag{3}$$

$$H'(t) = \gamma \mathrm{e}^{-\mathrm{i}\omega t} d^\dagger_\uparrow d_\downarrow + \mathrm{h.c.}, \tag{4}$$

$$H_\mathrm{T} = \sum_{kk'\sigma} t_\mathrm{LR}(C^\dagger_{\mathrm{L}k\sigma} C_{\mathrm{R}k'\sigma} + C^\dagger_{\mathrm{R}k'\sigma} C_{\mathrm{L}k\sigma}) + \sum_{k\sigma} (t_{\mathrm{L}d} C^\dagger_{\mathrm{L}k\sigma} d_\sigma + t_{\mathrm{R}d} \mathrm{e}^{\mathrm{i}\sigma\varphi} C^\dagger_{\mathrm{R}k\sigma} d_\sigma) + \mathrm{h.c.}, \tag{5}$$

where $H_\alpha (\alpha = \mathrm{L, R})$ describes the left and right normal leads with $\epsilon_{\alpha k}$ the single spin-degenerate electron energy in the α lead; $C^\dagger_{\alpha k\sigma} (C_{\alpha k\sigma})$ is the creation (annihilation) operator; H_D models the QD which consists of a single energy level ε_d and a Zeeman splitting $\sigma\mu_\mathrm{B} B_z$ under the z component of the RMF. $d^\dagger_\sigma (d_\sigma)$ denotes the creation (annihilation) operator for the electron with spin σ in the QD; $H'(t)$ is the transverse component of the RMF with $\gamma = \mu_\mathrm{B} B_0 \sin\theta$; H_T denotes the tunneling part of the Hamiltonian where t_LR, $t_{\mathrm{L}d}$, $t_{\mathrm{R}d}$ are all the tunneling coefficients and assumed to be real for easy derivation; the extra phase $\mathrm{i}\sigma\varphi$ arises from the existence of the RSOC.

We proceed to work out the pumped current in the device described above and stress again that there is no bias applied on the system. The electronic current per spin channel flowing from the left lead to the ring can be obtained from the time evolution of the occupation number for electrons with the corresponding spin in the left lead, i.e.

$$I_{\mathrm{L}\sigma}(t) = -e\langle \mathrm{d}N_{\mathrm{L}\sigma}/\mathrm{d}t \rangle = \frac{\mathrm{i}e}{\hbar} \langle [N_{\mathrm{L}\sigma}, H] \rangle, \quad N_{\mathrm{L}\sigma} = \sum_k C^\dagger_{\mathrm{L}k\sigma} C_{\mathrm{L}k\sigma}, \tag{6}$$

where $[\cdots]$ and $\langle\cdots\rangle$ denote operator commutation and thermal average, respectively. Using the Keldysh Green's function method, equation (6) can be expressed as

$$I_{L\sigma}(t) = \frac{e}{\hbar}[t_{Ld}G^{<}_{d\sigma,L\sigma}(t,t) + t_{LR}G^{<}_{R\sigma,L\sigma}(t,t) + \text{c.c.}], \quad (7)$$

where $G^{<}(t,t')$ is the lesser Green's functions defined as [25]

$$G^{<}_{\alpha\sigma,\alpha'\sigma'}(t,t') = \mathrm{i}\left\langle \sum_{k'} C^{\dagger}_{\alpha'k'\sigma'}(t') \sum_{k} C_{\alpha k\sigma}(t)\right\rangle, \quad (8)$$

$$G^{<}_{\alpha\sigma,d\sigma'}(t,t') = \mathrm{i}\left\langle d^{\dagger}_{\sigma'}(t') \sum_{k} C_{\alpha k\sigma}(t)\right\rangle, \quad (9)$$

$$G^{<}_{d\sigma,d\sigma'}(t,t') = \mathrm{i}\langle d^{\dagger}_{\sigma'}(t') d_{\sigma}(t)\rangle. \quad (10)$$

By performing a double-time Fourier transform, the time-averaged current with respect to $I_{L\sigma}(t)$ is given by [29]

$$I_{L\sigma} = \frac{2e}{\hbar}\frac{1}{2N\tau}\,\mathrm{Re}\int \frac{\mathrm{d}E}{2\pi}[t_{Ld}G^{<}_{d\sigma,L\sigma}(E,E) + t_{LR}G^{<}_{R\sigma,L\sigma}(E,E)], \quad (11)$$

where $G^{<}(E,E)$ is the Fourier transform of $G^{<}(t,t)$, $\tau = 2\pi/\omega$ is the rotating period of the RMF and $N \to \infty$. To solve the lesser Green's function, we first consider the contour ordered Green's function in Keldysh space defined as [30–32]

$$G^{c}(t,t') = \begin{pmatrix} G^{t}(t,t') & G^{<}(t,t') \\ G^{>}(t,t') & G^{\bar{t}}(t,t') \end{pmatrix}. \quad (12)$$

where $G^{>}$, G^{t}, and $G^{\bar{t}}$ are the greater, time ordered, and antitime ordered Green's functions, respectively,

$$G^{>}_{\alpha\sigma,\alpha'\sigma'}(t,t') = -\mathrm{i}\left\langle \sum_{k} C_{\alpha k\sigma}(t) \sum_{k'} C^{\dagger}_{\alpha'k'\sigma'}(t')\right\rangle, \quad (13)$$

$$G^{>}_{\alpha\sigma,d\sigma'}(t,t') = -\mathrm{i}\left\langle \sum_{k} C_{\alpha k\sigma}(t) d^{\dagger}_{\sigma'}(t')\right\rangle, \quad (14)$$

$$G^{>}_{d\sigma,d\sigma'}(t,t') = -\mathrm{i}\langle d_{\sigma}(t) d^{\dagger}_{\sigma'}(t')\rangle, \quad (15)$$

$$G^{t}(t,t') = \theta(t-t')G^{>}(t,t') + \theta(t'-t)G^{<}(t,t'), \quad (16)$$

and

$$G^{\bar{t}}(t,t') = \theta(t'-t)G^{>}(t,t') + \theta(t-t')G^{<}(t,t'). \quad (17)$$

The four component Green's functions of G^{k} are not completely independent. They are related by

$$G^{t} = G^{<} + G^{r}, \quad (18)$$

$$G^{\bar{t}} = G^{<} - G^{a}, \quad (19)$$

and

$$G^{>} = G^{t} - G^{a}. \quad (20)$$

Here $G^{r(a)}$ is the usual retarded (advanced) Green's function and its definition is

$$G^{r(a)}_{\alpha\sigma,\alpha'\sigma'}(t,t') = \mp\mathrm{i}\theta(\pm t \mp t')\left\langle\left[\sum_{k} C_{\alpha k\sigma}(t), \sum_{k'} C^{\dagger}_{\alpha'k'\sigma'}(t')\right]\right\rangle, \quad (21)$$

and $G^{r(a)}_{\alpha\sigma,d\sigma'}(t,t')$, $G^{r(a)}_{d\sigma,d\sigma'}(t,t')$ are defined in the same manner.

In a time-dependent problem, it is convenient to adopt the perturbation theory to calculate the contour ordered Green's function G^{c} because it fulfils the Dyson equation

$$G^{c}(t,t') = G^{c0}(t,t') + \int \mathrm{d}t_1 G^{c0}(t,t_1)V^{c}(t_1)G^{c0}(t_1,t') + \cdots, \quad (22)$$

with G^{c0} being the contour ordered Green's function when the time-dependent potential V^{c} is zero. With the Fourier transform, we can obtain the second order of $(v^{c})^{2}$ with the following equation

$$G^{c}(E,E') = G^{c0}(E,E') + \Delta^{(1)}G^{c}(E,E') + \Delta^{(2)}G^{c}(E,E') = 2\pi G^{c0}(E)\delta(E-E') + G^{c0}(E)V^{c}(E-E')G^{c0}(E') + \int \frac{\mathrm{d}E_1}{2\pi} G^{c0}(E)V^{c}(E-E_1)G^{c0}(E_1) \times V^{c}(E_1-E')G^{c0}(E'), \quad (23)$$

where $\Delta^{(1)}G^{c}(E,E')$ and $\Delta^{(2)}G^{c}(E,E')$ are the first-order and second-order corrections to the unperturbed contour ordered Green's function $G^{c0}(E,E')$, respectively. In the perturbation formula above, we treat the time-dependent component of the RMF $H'(t)$ as the perturbation and in the Keldysh space the perturbation potential $V^{c}(E)$ is defined as

$$V^{c}(E) = \begin{pmatrix} V(E) & 0 \\ 0 & -V(E) \end{pmatrix}, \quad (24)$$

where $V(E)$ in spin space is given by $V_{d\sigma,d\bar{\sigma}}(E) = 2\pi\gamma\delta(E - \sigma\omega)(\sigma = \pm;\uparrow,\downarrow$ and $\bar{\sigma} = -\sigma)$, which can be directly obtained from the Fourier transform of $H'(t)$. In the later calculations, the unperturbed contour ordered Green's function G^{c0} can be neglected because without the perturbation, the system is in the equilibrium state and G^{c0} can only result in a persistent spin current circulating in the ring [33] whose magnitude is negligibly small. Moreover, through our derivation, the first-order correction $\Delta^{(1)}G^{c}$ in equation (23) is found not to contribute to the generation of any current [31]. Thus when evaluating $G^{<}$ that is directly related to the spin-resolved current $I_{L\sigma}$ in equation (11), we only need keep its second-order correction, and we can get

$$\Delta^{(2)}G^{<} \sim G^{r0}V^{c}G^{r0}V^{c}G^{<0} + G^{r0}V^{c}G^{<0}V^{c}G^{a0} + G^{<0}V^{c}G^{a0}V^{c}G^{a0}, \quad (25)$$

where G^{r0}, G^{a0}, and $G^{<0}$ are, respectively, the retarded, advanced, and lesser Green's functions in the equilibrium state without the transverse RMF and they satisfy the well-known relation $G^{<0} = f(E)(G^{a0} - G^{r0})$ with the Fermi distribution function $f(E)$. With these preparations above, we obtain

$$I_{L\sigma} = I_{L_1\sigma} + I_{L_2\sigma}, \quad (26)$$

$$I_{L_1\sigma} = \frac{2e\gamma^2}{\hbar}\,\mathrm{Re}\int \frac{\mathrm{d}E}{2\pi} t_{Ld}[f(E) - f(E-\sigma\omega)] \times \{G^{r0}_{d\sigma,d\sigma}(E)[G^{r0}_{d\bar{\sigma},d\bar{\sigma}}(E-\sigma\omega) - G^{a0}_{d\bar{\sigma},d\bar{\sigma}}(E-\sigma\omega)]G^{a0}_{d\sigma,L\sigma}(E)\}, \quad (27)$$

and

$$I_{L2\sigma} = \frac{2e\gamma^2}{\hbar}\,\mathrm{Re}\int \frac{dE}{2\pi} t_{LR}[f(E) - f(E-\sigma\omega)] \times \{G^{r0}_{R\sigma,d\sigma}(E)[G^{r0}_{d\bar\sigma,d\bar\sigma}(E-\sigma\omega) - G^{a0}_{d\bar\sigma,d\bar\sigma}(E-\sigma\omega)]G^{a0}_{d\sigma,L\sigma}(E)\}, \quad (28)$$

where $I_{L1\sigma}$ and $I_{L2\sigma}$ denote the spin-resolved currents flowing through the arm with the dot (upper arm) and the arm without the dot (lower arm), respectively.

To solve $G^{r(a)0}$, we use the Dyson equation

$$G^{r(a)0}_\sigma = g^{r(a)}_\sigma + g^{r(a)}_\sigma \Sigma^{r(a)}_\sigma G^{r(a)0}_\sigma, \quad (29)$$

where the equilibrium retarded (advanced) Green's function $G^{r(a)0}_\sigma$ in the local basis is a 3×3 matrix

$$G^{r(a)0}_\sigma = \begin{bmatrix} G^{r(a)0}_{L\sigma,L\sigma} & G^{r(a)0}_{L\sigma,R\sigma} & G^{r(a)0}_{L\sigma,d\sigma} \\ G^{r(a)0}_{R\sigma,L\sigma} & G^{r(a)0}_{R\sigma,R\sigma} & G^{r(a)0}_{R\sigma,d\sigma} \\ G^{r(a)0}_{d\sigma,L\sigma} & G^{r(a)0}_{d\sigma,R\sigma} & G^{r(a)0}_{d\sigma,d\sigma} \end{bmatrix}. \quad (30)$$

By the equation of motion method [34] and with the wide-band approximation [35], the Green's function of the isolated system (i.e. when $t_{LR} = t_{Ld} = t_{Rd} = 0$) $g^{r(a)}_\sigma$ is given by

$$g^{r0}_\sigma(E) = \begin{bmatrix} -i\pi\rho & 0 & 0 \\ 0 & -i\pi\rho & 0 \\ 0 & 0 & \frac{1}{E-\varepsilon_d-\sigma\mu_B B_z + i0^+} \end{bmatrix}, \quad (31)$$

and

$$g^{a0}_\sigma(E) = [g^{r0}_\sigma(E)]^\dagger, \quad (32)$$

with ρ denoting the density of states of the leads at the Fermi level. For our present case, the self-energy is from the coupling between the two leads and the ring,

$$\Sigma^r_\sigma(E) = \begin{bmatrix} 0 & t_{LR} & t_{Ld} \\ t^*_{LR} & 0 & t_{Rd}e^{i\sigma\varphi} \\ t^*_{Ld} & t^*_{Rd}e^{-i\sigma\varphi} & 0 \end{bmatrix}. \quad (33)$$

With equations (27)–(33) and some direct algebra, we can obtain the charge and spin currents flowing though the left lead as

$$I_{L1} = I_{L1\uparrow} + I_{L1\downarrow} = 0, \quad (34)$$

$$J_{L1} = I_{L1\uparrow} - I_{L1\downarrow} = \frac{8e\gamma^2}{h} t^2 t^2_{Ld}(t^2_{Ld} + t^2_{Rd})(t^2 + t^2_{LR}) \times \int dE \frac{f(E) - f(E-\omega)}{X_\uparrow(E)X_\downarrow(E-\omega)}, \quad (35)$$

$$I_{L2} = I_{L2\uparrow} + I_{L2\downarrow} = \frac{8e\gamma^2 \sin\varphi}{h} t t_{LR} t_{Ld} t_{Rd}(t^2_{Ld} + t^2_{Rd})(t^4 - t^4_{LR}) \times \int dE \frac{f(E) - f(E-\omega)}{X_\uparrow(E)X_\downarrow(E-\omega)}, \quad (36)$$

$$J_{L2} = I_{L2\uparrow} - I_{L2\downarrow} = \frac{8e\gamma^2}{h} t^2 t^2_{LR}(t^2 + t^2_{LR})(t^4_{Rd} - t^4_{Ld}) \times \int dE \frac{f(E) - f(E-\omega)}{X_\uparrow(E)X_\downarrow(E-\omega)}, \quad (37)$$

with $t = 1/\pi\rho$ and $X_\sigma(E) = [(E - \varepsilon_d - \sigma\mu_B B_z)(t^2 + t^2_{LR}) + 2t_{LR}t_{Ld}t_{Rd}\cos\varphi]^2 + t^2(t^2_{Ld} + t^2_{Rd})^2$. Here $I_{L1(2)}$ and $J_{L1(2)}$ represent the tunneling charge and spin current flowing from the left lead to the upper (lower) arm, and subsequently the total charge and spin currents are given by $I_L = I_{L1} + I_{L2}$, $J_L = J_{L1} + J_{L2}$. Equations (34)–(37) form the main results of this paper, which indicate clearly that a nonzero spin and charge current can flow in the one-dot Rashba ring device without external bias.

The generated currents have some conspicuous characteristics. Firstly, only a pure spin current J_{L1} can flow in the upper ring arm (see equations (34) and (35)), which cannot contribute to the charge current but only to the spin current in the left lead. This pure spin current J_{L1} is driven by the RMF, and it depends on the strength of the transverse component of the RMF γ as well as the rotating frequency ω. The RSOC in the ring can only affect the magnitude of the pumped spin current J_{L1} through the quantity $X_\sigma(E)$. Secondly, in the lower arm without QD, a nonzero charge current I_{L2} emerges accompanied by a spin current J_{L2} so that I_{L2} is generally spin-polarized. In the same way as J_{L1}, the spin current J_{L2} is also pumped by the RMF and can be modulated by device parameters such as the RMF strength γ, the rotating frequency ω, the tunneling coefficients t_{Ld}, t_{Rd}, and t_{LR}. Especially, if $t_{Ld} = t_{Rd}$, the spin current J_{L2} will vanish, which indicates that the current flowing in the lower arm comes from the interference effect since at this symmetric case the spin current is absent, therefore, the spin-polarization of I_{L2} can be modulated by device parameters and can even be unpolarized or fully-polarized. Different from the current flowing in the upper arm where the pure spin current is driven by the RMF, the RSOC is decisive to the induced charge current I_{L2}, i.e. when there is no RSOC or $\varphi = n\pi$ with n an integer, the charge current will disappear, since the charge current I_{L2} is the result of the interplay between the spin-pump effect driven by RMF and the quantum interference effect in the ring induced by the spin precession phase φ, thus the tunneling coefficients t_{Ld}, t_{Rd}, and t_{LR} are all prerequisites for the generation of the charge current as shown in equation (36). Finally, since $I_L = I_{L1} + I_{L2}$, $J_L = J_{L1} + J_{L2}$, the charge current in the left lead is only contributed by the one in the lower arm (without a QD), i.e. $I_L = I_{L2}$, and must be spin-polarized. Here, it should be noted that though $I_L = I_{L2}$, spin-polarized charge currents I_L and I_{L2} have different spin-polarization degrees, because the inequality $J_L \neq J_{L2}$ always holds according to equation (37) and the relation $J_L = J_{L1} + J_{L2}$.

3. Numerical calculations and discussions

In this section, we present our numerical results of the pumped current based on equations (35)–(37). In all numerical calculations the Fermi energy of the system is taken as the energy reference, $E_F = 0$, and $t = 1/\pi\rho = 1$ is taken as the energy unit, the temperature is zero $T = 0$. We first present the charge current I_L as a function of the single energy level in the QD ε_d for different spin precession phases φ, since ε_d can be easily controlled by a gate voltage in reality. As is shown in figure 2, the charge current varies sensitively with the dot level and has a sizable value at its resonant peak, which

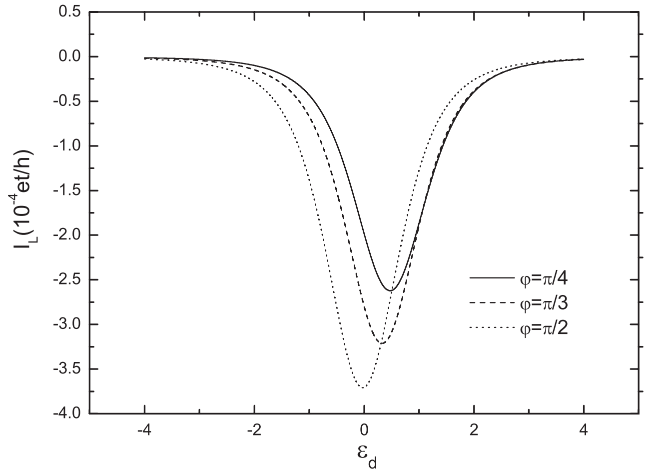

Figure 2. The induced charge current $I_{\rm L}$ as a function of the dot level ε_d for different spin precession phases. The other parameters are $\mu_{\rm B}B_z = 0.1$, $\gamma = 0.1$, $\omega = 0.05$, $t_{\rm LR} = 0.6$, $t_{\rm L\it d} = 0.8$, and $t_{\rm R\it d} = 1$.

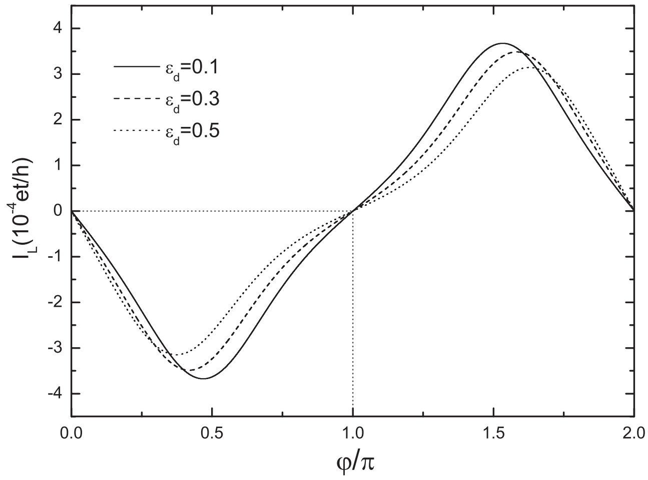

Figure 3. The induced charge current $I_{\rm L}$ as a function of the spin precession phase φ for different dot levels. The other parameters are the same as those in figure 2.

is determined by ε_d and φ. The RSOC can affect both the magnitude of $I_{\rm L}$ and its resonant energy, which is clearly shown in the quantity X_E in equations (35)–(37), and obviously the resonant energy is at $E_{\rm F} = \varepsilon_d$ with $\varphi = \pi/2$.

The charge current $I_{\rm L}$ versus the precession phase φ is plotted in figure 3. Different from that shown in figure 2, both the magnitude and the sign of the charge current can vary with φ. As can be seen, it is a sinusoid-like function of the spin precession phase φ and oscillates with period 2π. When $\varphi = 0$, i.e. the RSOC is absent, the charge current disappears. This is because the charge current results from the transfer from the spin current pumped by the RMF due to the quantum interference effect of the RSOC induced spin precession phase. According to equation (36), at $\varphi = n\pi (n = \pm1, \pm2, \ldots)$, the quantum interference in the ring reaches its destructive points and thus the charge current also disappears, as shown in figure 3. Since the RSOC strength can be tuned by a gate voltage in experiments [27, 28], we can modulate the charge current in its magnitude and direction electrically.

In figure 4, the charge current $I_{\rm L}$ is depicted as a function of the rotating frequency ω. As is known, the mono-parametric charge pump cannot survive in the adiabatic regime ($\omega \sim 0$) and has a quadratic dependence on ω in the finite pumping frequency regime when the spatial inversion symmetry of the system is broken. In our device, the current direction can be modulated by φ, which is related to the spatial configuration of the system. When we carry the inversion operation on the system, $\varphi \to -\varphi$ and $I_{\rm L} \to -I_{\rm L}$ although the device itself remains unchanged. As shown in figure 4, the charge current $I_{\rm L}$ is linear with ω at low frequency, which reflects that the charge current comes from the interplay between the spin-pump effect and the quantum interference effect since the pumped pure spin current by the RMF is linearly dependent on ω at low frequency. Therefore, the generated charge current in our device does not vanish at the adiabatic limit although there is only one pumping source.

We proceed to investigate the spin-polarization degrees of $I_{\rm L2}$ and $I_{\rm L}$, which are defined as $\eta_2 = J_{\rm L2}/I_{\rm L2}$ and $\eta = J_{\rm L}/I_{\rm L}$ respectively. Figure 5 shows the spin-polarization degree η_2 of the current flowing in the lower arm of the ring. As is shown, η_2 is dependent on the tunneling coefficient $t_{\rm L\it d}$ and can even reverse its sign. Of particular interest is the case of $\eta_2 = \pm1$, i.e. the charge current is fully spin-polarized and only one spin species current can flow in the lower arm of the ring, i.e. $I_{\rm L2\uparrow} \neq 0$ and $I_{\rm L2\downarrow} = 0$ as $\eta_2 = 1$ and vice versa as $\eta_2 = -1$. As $t_{\rm L\it d} = t_{\rm R\it d} = 1$ irrespective of whatever other parameters are, the spin-polarization degree $\eta_2 = 0$ and the charge current is not spin-polarized, as stated early, the spin current of $J_{\rm L2}$ vanishes at the symmetric coupling $t_{\rm L\it d} = t_{\rm R\it d}$ due to the quantum interference effect. From figure 5, it can be seen that the spin-polarization in the lower arm can be arbitrarily controlled, which is useful in spintronics to manipulate the spin efficiently.

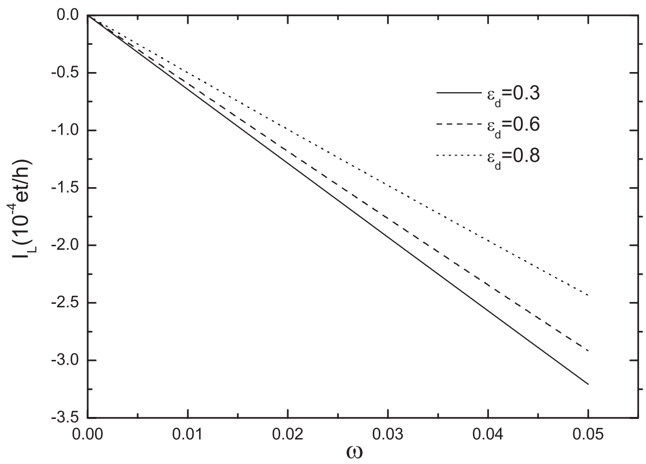

Figure 4. The dependence of current $I_{\rm L}$ on the rotating frequency ω for different dot levels. The spin precession phase is taken as $\varphi = \pi/3$. The other parameters are the same as those in figure 2.

The spin-polarization degree of the total current in the left lead η is presented in figure 6 as a function of the spin precession phase φ. It shows that with an appropriate choice of the other parameters, η can vary with φ from -1 to 0 and

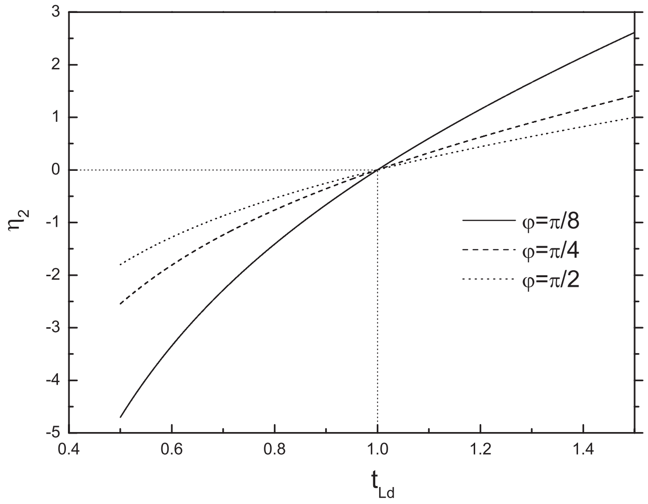

Figure 5. The spin-polarization degree η_2 versus the tunneling coefficient t_{Ld} for different spin precession phases. The other parameters are $\varepsilon_d = 0.01$, $\mu_B B_z = 0.1$, $\gamma = 0.1$, $\omega = 0.05$, $t_{LR} = 1.5$, and $t_{Rd} = 1$.

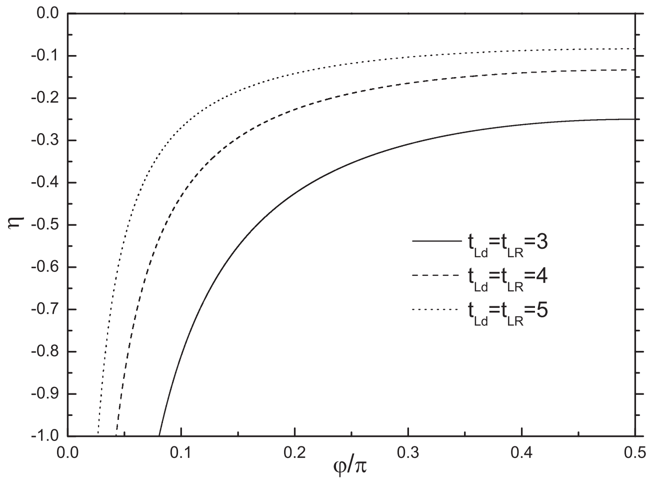

Figure 6. The spin-polarization degree η dependence on the spin precession phase for different tunneling coefficients t_{Ld} and t_{LR}. The other parameters are $\varepsilon_d = 0.1$, $\mu_B B_z = 0.1$, $\gamma = 0.1$, $\omega = 0.05$, and $t_{Rd} = 1$.

the zeroth value of η can only be approached but cannot be reached, because in the presence of the RMF the spin current J_{L1} always exists and contributes to the spin current J_L so that the charge current I_L can never be unpolarized. Hence, the fully spin-polarized charge current can also flow in the lead besides I_{L2} discussed above. It is noted that the value of η can vary between 0 and 1 when the phase φ in figure 6 shifts with π. In other words, the spin-polarization of the charge current in the left lead can also be almost arbitrarily modulated between -1 and $+1$ (except 0) by tuning the RSOC strength with the other parameters fixed.

For a uniform RSOC ring [36], the AC phase φ is estimated to be 7.4π for a 0.3 μm ring radius with the RSOC constant $\alpha \sim 1.05 \times 10^{-11}$ eV m, which is sufficient for modulating the generated spin-polarized current. Such a size is also enough to maintain the quantum spin-interference effect in the ring. Hence, the proposed device can work well so long as it is kept clean and at an appropriate temperature.

4. Conclusion

In summary, we have studied the spin-polarized transport through a two-terminal one-dot Rashba ring device driven by an RMF. Due to the coexistence of the pure spin current pumped by the RMF and the spin-dependent quantum interference effect in the ring from the AC phase, a spin-polarized charge current, which is proportional to the rotating frequency of the RMF, can be generated in the lead and the arm without the QD. By tuning the device parameters such as the dot energy level, the RSOC strength, and the tunneling strength, the spin-polarized charge current can be easily modulated in magnitude, direction, and even its spin-polarization degree. The proposed device may act as an efficient generator of a tunable spin-polarized current with arbitrary spin-polarization degree in spintronics.

Acknowledgments

The work was supported by the General Research Fund of the Research Grants Council of Hong Kong SAR, China (Project No. CityU 100308/08P); JW acknowledges support from NSFC (10704016) and NBRPC (2009CB929504); YHY acknowledges support by NSF of JiangSu (BK2007100).

References

[1] Datta S and Das B 1999 *Appl. Phys. Lett.* **75** 695
[2] Tsukagoshi K, Alphenaar B W and Ago H 1999 *Nature* **401** 572
[3] Wolf S A, Awschalom D D, Buhrman R A, Daughton J M, von Molnár S, Roukes M L, Chtchelkanova A Y and Treger D M 2001 *Science* **294** 1488
[4] Koga T, Nitta J and Takayanagi H 2002 *Phys. Rev. Lett.* **88** 126601
[5] Schliemann J, Carlos Egues J and Loss D 2003 *Phys. Rev. Lett.* **90** 146801
[6] Kikkawa J M, Smorchkova I P, Samarth N and Awschalom D D 1997 *Science* **277** 1284
Kikkawa J M and Awschalom D D 1999 *Nature* **397** 139
[7] Murakami J S, Nagaosa N and Zhang S C 2003 *Science* **301** 1348
[8] Sinova J, Culcer D, Niu Q, Sinitsyn N A, Jungwirth T and MacDonald A H 2004 *Phys. Rev. Lett.* **92** 126603
[9] Johnson M and Silsbee R H 1985 *Phys. Rev. Lett.* **55** 1790
[10] Fiederling R, Kein M, Reuscher G, Ossau W, Schmidt G, Waag A and Molenkamp L W 1999 *Nature* **402** 787
[11] Ohno Y, Young D K, Beschoten B, Matsukura F, Ohno H and Awschalom D D 1999 *Nature* **402** 790
[12] Jedema F J, Filip A T and van Wees B J 2001 *Nature* **410** 345
[13] Wang B G, Wang J and Guo H 2003 *Phys. Rev.* B **67** 092408
[14] Zhang P, Xue Q K and Xie X C 2003 *Phys. Rev. Lett.* **91** 196602
[15] Brataas A, Tserkovnyak Y, Bauer G E W and Halperin B I 2002 *Phys. Rev.* B **66** 060404(R)
[16] Hattori K 2007 *Phys. Rev.* B **75** 205302
[17] Li C, Yu Y, Wei Y and Wang J 2007 *Phys. Rev.* B **75** 035312
[18] Zhou Y Q, Wang R Q, Sheng L, Wang B G and Xing D Y 2008 *Phys. Rev.* B **78** 155327

[19] Frustaglia D, Hentschel M and Richter K 2001 *Phys. Rev. Lett.* **87** 256602
[20] Hentschel M, Schomerus H, Frustaglia D, Hentschel M and Richter K 2004 *Phys. Rev.* B **69** 155326
[21] Wu M W, Zhou J and Shi Q W 2004 *Appl. Phys. Lett.* **85** 1012
[22] Ionicioiu R and D'Amico I 2003 *Phys. Rev.* B **67** 041307(R)
[23] Citro R and Romeo F 2008 *Phys. Rev.* B **77** 193309
[24] Chen K W and Chang C R 2008 *Phys. Rev.* B **78** 235319
[25] Sun Q F, Wang J and Guo H 2005 *Phys. Rev.* B **71** 165310
Sun Q F and Xie X C 2006 *Phys. Rev.* B **73** 235301
[26] Sun Q F and Xie X C 2005 *Phys. Rev.* B **71** 155321
[27] Governale M, Taddei F and Fazio R 2003 *Phys. Rev.* B **68** 155324
[28] Mal'shukov A G, Tang C S, Chu C S and Chao K A 2003 *Phys. Rev.* B **68** 233307
[29] Wang B G, Wang J and Guo H 2003 *Phys. Rev.* B **68** 155326
[30] Keldysh L V 1965 *Sov. Phys.—JETP* **20** 1018
[31] Wang J and Chan K S 2008 *J. Phys.: Condens. Matter* **20** 055227
[32] Mahan G D 1993 *Many Particle Physics* (New York: Plenum)
[33] Sun Q F, Xie X C and Wang J 2007 *Phys. Rev. Lett.* **98** 196801
[34] Sun Q F, Xie X C and Wang J 2008 *Phys. Rev.* B **77** 035327
[35] Jauho A P, Wingreen N S and Meir Y 1994 *Phys. Rev.* B **50** 5228
[36] Nitta J, Meijer F E and Takayanagi H 1999 *Appl. Phys. Lett.* **75** 695