All-electrical measurement of relative strength of spin–orbit interactions in interacting quantum wires

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New Journal of Physics 14 (2012) 013016 (10pp)
Received 2 May 2011
Published 11 January 2012
Online at http://www.njp.org/
doi:10.1088/1367-2630/14/1/013016

Abstract. We present a theoretical analysis of the ac transport property of quantum wires with Coulomb interaction included using the Luttinger liquid theory. We make experimentally testable predictions of a sensitive dependence of the ac conductance on the wire orientation. The ac conductance displays oscillating behavior as the frequency of the external voltage increases. The oscillation period provides us with a new way of measuring the relative strength of the Rashba and Dresselhaus spin–orbit interactions.

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1. Introduction

Generating spin-polarized electrons electrically in semiconductors is one of the central issues and the ultimate goal of the spintronics field ([1] and references therein). The spin–orbit interaction (SOI), a manifestation of special relativity, provides us with an efficient way of controlling the electron spin via its orbital motion ([2] and references therein). The SOI can be induced by breaking the structure inversion symmetry or the bulk crystal inversion symmetry, which gives rise to, respectively, the Rashba SOI (RSOI) and Dresselhaus SOI (DSOI) [3–5]. The absolute strength of the RSOI can be tuned electrically and is usually much larger than that of the DSOI. However, in thin quantum wells, the strength of the DSOI is comparable to that of the RSOI since the strength of the DSOI depends sensitively on the thickness of the quantum well. There are many proposals for all-electrical spin filtering and generation of spin-polarized current in semiconductors based on the Rashba and Dresselhaus SOIs. It is very useful to have precise control of the Dresselhaus and Rashba SOIs [6], as the interplay between them can lead to some remarkable properties, such as an extremely long spin lifetime, the anisotropic photogalvanic effect [7] and the persistent spin helix [8]. As a result, it is crucial to know accurately the strengths of the two SOIs in semiconductors. Some experimental works have already measured the relative strengths of the RSOI and DSOI utilizing the photogalvanic effect [7] and the time-resolved Faraday rotation [9].

Very recently, an extremely long spin lifetime was found in semiconductor quantum wires where the effective magnetic field induced by the SOIs is fixed along a specific crystallographic direction [10]. A very recent paper proposed a new scheme for detecting the relative strengths of the RSOI and DSOI utilizing the orientation dependence of the dc conductance in quasi-one-dimensional (Q1D) semiconductor quantum wires [11]. However, the effect of the Coulomb interaction on the transport property has not been addressed and this paper presents a study of this effect. Since the Coulomb interaction has a pronounced effect in Q1D systems, the conventional Fermi liquid (FL) theory breaks down, and we need to use the Luttinger liquid (LL) theory in our study [12]. LL behavior was demonstrated experimentally in many Q1D systems, e.g. the spin–charge separation [13, 14], and some recent studies have investigated quantum wire systems including the effect of the RSOI [15–17].

In this paper, based on the LL theory, we propose a new scheme for measuring the relative strengths of the RSOI and DSOI using the sensitive dependence of the ac conductance on wire orientation. The wire orientation dependence of the dc conductance is smeared out when the quantum wire is connected perfectly to the FL leads. However, the ac conductance still exhibits a strong orientation dependence and interesting oscillation caused by the formation of Fabry–Pérot (FP) modes of the collective excitations in the central wire. This orientation-dependent behavior of the ac conductance depends sensitively on the strengths of the SOIs. From the orientation dependence, one can extract the relative strength of the RSOI and DSOI, i.e. the R/D ratio.

2. Model and theory

First we consider a two-probe configuration: a quantum wire under a top gate connected to two long leads in order to simulate two reservoirs with fixed Fermi energies (see figure 1(a)). The recent progress in the semiconductor technique makes it possible to fabricate quantum wires with long coherent length; one can investigate ballistic transport in such quantum wires. The
Figure 1. (a) Schematic of a semiconductor quantum wire fabricated using the split gate technique along the crystallographic direction $\theta$ with respect to the [100] axis in the crystallographic planes (001). $n$ is the unit vector along the normal direction of the crystallographic plane. The RSOI only exists in the region below the electric gate, while the DSOI exists everywhere in the sample. (b) The single-electron energy spectrum of 2DEG in the presence of the RSOI and DSOI. (c) The constant energy surfaces and the corresponding spin orientations.

Hamiltonian of the noninteracting electrons in the Q1D quantum wire oriented along different crystallographic directions in the (001) plane is

$$H_0 = \frac{\hbar^2 k_x^2}{2m^*} - \alpha \sigma_y k_x + \beta [\cos(2\theta)\sigma_z k_x - \sin(2\theta)\sigma_y k_x], \quad (1)$$

where $\theta$ is the angle between the quantum wire and the [100] axis, $m^*$ is the electron effective mass, $\sigma_i$ ($i = x, y, z$) are the Pauli matrices and $\alpha$ and $\beta$ are the strengths of the RSOI and DSOI, respectively. Consider the case when electrons occupy only the lowest subband of the quantum wire, and high above the bottom of the subband, the linearized noninteracting electron Hamiltonian of the quantum wire with both RSOI and DSOI is $H_0 = -i\hbar \int \sum_{\gamma, s} v_{\gamma s} \psi_{\gamma s}^\dagger \psi_{\gamma s} dx$, where the operators $\psi_{\gamma s} (\gamma = -1(L), 1(R); s = -1(\downarrow), 1(\uparrow))$ annihilate spin-down ($\downarrow$) or spin-up ($\uparrow$) electrons near the left (L) and right (R) Fermi points, respectively. $v_{\gamma s} = \gamma v_F - s \delta v$ are the four different Fermi velocities, where $v_F$ is the bare Fermi velocity of the right and left moving noninteracting electrons, and $\delta v = \sqrt{\alpha^2 + \beta^2 + 2\alpha\beta \sin 2\theta / \hbar}$. Note that the RSOI and DSOI split the electron subbands and make the electron Fermi velocities spin and orientation dependent. The model Hamiltonian (1) is valid when the lowest spin-up and spin-down subbands are occupied and the energy spacing between the lowest and the first excited subband is much larger than the Fermi energy, i.e. the narrow quantum wire case. When higher subbands are involved, anticrossing occurs between the lowest subband and the first excited subband and changes the spin orientation of electrons; therefore the Hamiltonian will be different [18]. In this paper, we focus only on the one-channel ballistic wires, i.e. the occupied lowest subband (strictly speaking,
we consider the case when the Fermi energy is located below the anticrossing point. The Coulomb interaction is given by $H_{\text{int}} = \frac{1}{2} \int \int dxdy \psi_{\gamma s}(x) \psi_{\gamma' s}(y) V_{\gamma' s}(x-y) \psi_{\gamma' s}(y) \psi_{\gamma s}(x)$. We can bosonize $H = H_0 + H_{\text{int}}$ and obtain

$$H = \frac{\hbar}{2} \int dx \left[ \frac{v_\rho}{K_\rho} (\partial_x \vartheta_\rho)^2 + v_\rho K_\rho \left( \frac{\Pi_\rho}{\hbar} \right)^2 \right] + \frac{\hbar}{2} \int dx \left[ \frac{v_\sigma}{K_\sigma} (\partial_x \vartheta_\sigma)^2 + v_\sigma K_\sigma \left( \frac{\Pi_\sigma}{\hbar} \right)^2 \right]$$

$$+ \hbar \int dx \delta v \left[ \left( \frac{\Pi_\sigma}{\hbar} \right) (\partial_x \vartheta_\sigma) + \left( \frac{\Pi_\rho}{\hbar} \right) (\partial_x \vartheta_\rho) \right],$$

(2)

where $\vartheta_\rho$ and $\vartheta_\sigma$ are the phase fields for the charge and spin degrees of freedom, respectively, and $\Pi_\rho$ and $\Pi_\sigma$ are the corresponding conjugate momenta. $v_{\rho,\sigma}$ are the propagating velocities of the charge and spin collective modes of the decoupled case $(\delta v = 0)$, $v_{\rho,\sigma} = v_F / K_{\rho,\sigma}$ and the parameter $K_{\rho/\sigma}$ is defined as $1/K_{\rho/\sigma}^2 = 1 \pm g$, where $g = 2V(q = 0) / \hbar \pi v_F$ with $V(q = 0)$ being the electron–electron interaction. In realistic samples, the interaction parameter $K$ varies spatially near the interface between the wire and the reservoirs. The spatial variation of $K$ is within a characteristic smoothing length $L_s$, and $L \gg L_s \gg \lambda_F$, where $\lambda_F$ is the electron Fermi wavelength, and $L$ is the length of wire. We can represent the variation of the parameter $K$ near the interface by a step function. The precise form of the variation of the parameter $K$ is not essential to the underlying physics, i.e. the FP modes in the wire, and our conclusion qualitatively. The Umklapp scattering process can be safely neglected because the Fermi energy in quantum wires formed in semiconductor heterostructures is far from the half-filled case. The electron–electron backscattering in a clean wire can also be neglected for a sufficiently long interacting region [15].

As shown schematically in figure 1, we consider a quantum wire with finite length, which is connected to the source and drain where the electron–electron interaction is negligible. The DSOI still exists in the source and drain because the wire is formed in a zincblende semiconductor heterostructure. The quantum wire is driven by a time-dependent source–drain voltage $H_{\text{ac}} = -\sqrt{\frac{2}{e}} \int dx E(x, t) \vartheta_\rho(x, t)$ is the Hamiltonian term that describes the effect of the time-dependent voltage on the charge excitation, where $-eE(x, t)$ can explicitly be replaced by the expression $\partial_x \mu = -\mu_L \delta(x) + \mu_R \delta(x - L)$ in which $\mu_{L/R}$ are the chemical potentials of the left and right leads, respectively [19]. The total Hamiltonian becomes $H = H_0 + H_{\text{int}} + H_{\text{ac}}$. The conductance of the quantum wire in an external electric field is

$$G(\omega) = \frac{1}{L^2} \int_0^L dx \int_0^L dx' \sigma_\rho(x, x', \omega),$$

(3)

where the ac conductivity $\sigma_\rho(x, x', \omega) = \frac{e^2}{2 m} \int_0^\infty d\tau e^{i\omega \tau} \langle T_\tau j_\rho(x, \tau) j_\rho(x', 0) \rangle |_{\omega \rightarrow \omega+i0}$. Mathematically, the conductivity $\sigma_\rho(0, L, \omega)$ is equivalent to the conductance for a Q1D quantum wire with finite length $L$. In terms of the charge continuity equation $i \partial \rho / \partial \tau + \partial j_\rho / \partial x = 0$ and the boson representation for $\rho$, one can obtain $\sigma_\rho(x, x', \omega) = \frac{4e^2}{\hbar} \omega_n G^{\rho \rho}_{\omega n}(x, x') |_{\omega \rightarrow \omega+i0}$, with $G^{\rho \rho}_{\omega n}(x, x')$ being the propagator of the boson field $\vartheta_\rho$, $G^{\rho \rho}_{\omega n}(x, x') = \frac{1}{\omega_n} \int_0^\infty d\tau e^{i\omega_n \tau} \langle T_\tau \delta \vartheta_\rho(x, \tau) \delta \vartheta_\rho(x', 0) \rangle$. The charge and spin degrees of freedom are coupled; thus the equation for $G^{\rho \rho}_{\omega n}(x, x')$ turns out to be coupled to the equation for $G^{\rho \rho}_{\omega n}(x, x')$:

$$\left[ -\frac{d}{dx} \left( f_\rho(x) \frac{d}{dx} \right) + \frac{\omega_n^2}{v_F} \right] G^{\rho \rho}_{\omega n}(x, x') + \frac{2\delta v \omega_n}{v_F} \frac{d}{dx} G^{\rho \rho}_{\omega n}(x, x') = \delta (x-x'),$$

(4)
Here, \( f_\pm(x) = (1 \pm g(x) - \delta v(x)^2/v_F^2) \), the parameter \( g(x) = g \) for \( 0 \leq x \leq L \) and zero otherwise, while \( \delta v(x) = \delta v \) for \( 0 \leq x \leq L \) and \( \beta/h \) otherwise. The propagator of the boson field \( G_{\omega n}^{\rho \sigma}(x, x') \) has a simple exponential form in all regions of

\[
G_{\omega n}^{\rho \sigma}(x, x') = A_{11} e^{-\frac{\pi}{L} x} + A_{22} e^{\frac{\pi}{L} x} + A_{33} e^{-\frac{\pi}{2L} x} + A_{44} e^{\frac{\pi}{2L} x},
\]

where \( A_{ii} \) (\( i = 1, 2, 3, 4 \)) can be deduced from the following boundary conditions: the functions \( G_{\omega n}^{\rho \sigma}(x, x') \), \( G_{\omega n}^{\rho \sigma}(x, x') \) and \( f_\pm(x) \frac{d}{dx} G_{\omega n}^{\rho \sigma}(x, x') \) are continuous at \( x = 0, x' \) and \( L \); the function \( f_\pm(x) \frac{d}{dx} G_{\omega n}^{\rho \sigma}(x, x') \) is continuous at \( x = 0, L \), but undergoes a jump at \( x = x' \), i.e.

\[
-\nu_F f_+(x) \frac{d}{dx} G_{\omega n}^{\rho \sigma}(x, x')|_{x=x'=0} = 1.
\]

In terms of the relation \( \sigma_\rho(x, x', \omega) = \frac{\delta v^2}{\pi} \omega_n G_{\omega n}^{\rho \sigma}(x, x')|_{\omega_n \rightarrow \omega+i0} \), we can finally obtain the ac conductivity of a quantum wire

\[
\sigma_\rho(x, x', \omega) = \frac{2e^2}{h} \left( A_1 e^{i \frac{\pi}{L} x} + A_2 e^{-i \frac{\pi}{L} x} + A_3 e^{i \frac{\pi}{2L} x} + A_4 e^{-i \frac{\pi}{2L} x} \right),
\]

where \( A_i = 2\omega_n A_{ii} \) is a function of \( x' \) and \( \omega \). The propagation velocities of the coupled charge and spin excitations are

\[
u^2_{1,2} = \frac{\nu^2_\rho + \nu^2_\sigma}{2} + \delta v^2 \pm \sqrt{\left( \frac{\nu^2_\rho - \nu^2_\sigma}{2} \right)^2 + 2\delta v^2 (\nu^2_\sigma + \nu^2_\rho)}.
\]

Note that the velocities \( u_{1,2} \) depend on the crystallographic orientation \( \theta \), which means that the SOIs couple the spin and charge excitations propagating independently in the quantum wire in the absence of the SOIs.

3. Numerical results and discussions

In the absence of the SOIs, i.e. \( \alpha = \beta = 0 \), the dc conductance \( G = 2K_L e^2/h \), where \( K_L \) is the interaction parameter in the leads, with \( K_L = 1 \) for the FL leads and \( K_L < 1 \) for the LL leads [20, 21]. In the presence of the SOIs, if the interacting quantum wire is infinitely long, the dc conductance depends sensitively on the crystallographic orientation \( \theta \) of the quantum wire, which can be referred to as the anisotropic transport behavior. But for a realistic sample, i.e. a narrow and clean quantum wire with finite length connected perfectly to the FL leads, the orientation dependence of the dc conductance disappears. From equations (3) and (7), one can see that the dc conductance is always equal to \( 2e^2/h \). It means that the SOIs do not affect the dc conductance of quantum wires connected to FL leads oriented in arbitrary directions, and this is the result of the contact resistance at the ends of the quantum wire. Therefore one actually cannot measure the strengths of the RSOI and DSOI utilizing the dc conductance, which is different from the noninteracting case [11].

As the dc conductance does not depend on orientation, it seems that it is impossible to measure the strengths of the RSOI and DSOI via the transport property of a quantum wire. So we propose a new scheme for measuring the relative strength of the RSOI and DSOI using the ac conductance. Different from dc conductance, which is entirely determined by the properties of the reservoirs, ac conductance is determined by transmission. For a quantum wire...
connected to the FL leads, the central quantum wire acts as an FP cavity in which electrons move back and forth and form resonant modes. The FP modes of charge and spin excitations are essentially caused by a mismatch of the strength of the electron–electron interaction and the SOIs in the quantum wire and the leads. These resonant modes, which are determined by the orientation-dependent Fermi velocity and the energy dispersion of the quantum wire, affect the electron transmission through the quantum wire. Therefore the ac conductance can still display a strong orientation dependence depending on the strengths of the SOIs. This significant variation of the real part of the ac conductance \( \text{Re}[G(\omega)] \) as a function of the wire orientation can be clearly seen in figure 2 for different strengths of the SOIs and the Coulomb interaction. Interestingly, the ac conductance oscillates with the changing crystallographic orientation of the quantum wire with a periodicity \( \pi \) for both the \( \alpha = \beta \) (see figure 2(a) or (b)) and \( \alpha \neq \beta \) cases (see figures 2(c)–(f)). This remarkable feature is caused by the different FP modes in the central wire determined by the orientation-dependent velocities of charge and spin collective excitations. Switching the strength of RSOI \( \alpha \) to \(-\alpha\), i.e. reversing the direction of the perpendicular electric field, will make the phase of the conductance oscillation shift relatively by \( \pi/2 \) (see the red and black lines in figure 2), which can be understood from the change of the velocity \( \delta v = \sqrt{\alpha^2 + \beta^2 + 2\alpha\beta \sin 2\theta/\hbar} \). The oscillation of the conductance is the interference effect of charge and spin plasmonic excitations. The FP modes of charge and spin excitations are essentially caused by the Andreev-like reflections and SOI-induced reflections at the contacts. The Andreev-like reflection of charge excitation at the contacts is because of the different interaction parameters \( K_\rho < 1 \) in the quantum wire and \( K_\rho = 1 \) in the FL leads). This reflection is momentum-conserving reflection and similar to the Andreev reflection at a

**Figure 2.** The ac conductance (in units of \( e^2/h \)) of a quantum wire with FL leads as a function of \( \theta \) with fixed \( \omega = 0.1 \times 2\pi v_F/L \): (a) for \( K_\rho = 0.4 \) where the black line is for \( \alpha = \beta = 0.2 \) and the red line for \( \alpha = -\beta = -0.2 \) (in units of […]
Figure 3. The ac conductance (in units of $e^2/h$) of a quantum wire with FL leads as a function of $K_\rho$ for fixed $\omega = 0.1 \times 2\pi v_F/L$, $\theta = 3\pi/4$. (a) $\alpha = \beta = 0.2$; (b) $\alpha = 0.3$, $\beta = 0.2$.

meta-superconductor contact. The SOIs induce an effective in-plane magnetic field pointing along the y-axis and polarize electron spins in the quantum wire and the leads. Note that in our model the RSOI is generated in the region below the top gate, i.e. in the quantum wire, and DSOI exists in both the quantum wire and the leads. The different SOIs in the quantum wire and the leads induce different effective magnetic fields and cause reflections of plasmonic excitations at the contacts. The effective magnetic fields induced by the RSOI and DSOI in the quantum wire cancel each other exactly (see equation (1)) at $\theta = 3\pi/4$ for $\alpha = \beta$ or $\theta = \pi/4$ for $\alpha = -\beta$. In this case, electrons in the wire become spin unpolarized suddenly, but electrons in the leads are polarized long the y-axis. Thus the conductance changes abruptly (see sharp spikes in figures 2(a) and (b)).

From the above results, we can find that the velocities $u_1$ and $u_2$ are important parameters in the ac conductance in the interacting wire. These velocities depend sensitively on the strengths of the SOIs and the Coulomb interaction, and lead to orientation-dependent FP modes in the interacting wire. In figure 3, we plot the real part of the ac conductance as a function of the strength of the Coulomb interaction $K_\rho$, for the quantum wire oriented along [1\(\bar{1}0\)] with two different relative strengths of the RSOI and DSOI. From the figure, one can see that tuning the strength of the Coulomb interaction $K_\rho$, i.e. the electron density, one can change the real part of the ac conductance significantly for the different strengths of the RSOI $\alpha$ and DSOI $\beta$. Note that the ‘slow’ collective excitation can approach zero, i.e. $u_2 = 0$ when $\delta v = v_\alpha = \sqrt{1 - g v_F}$. This feature corresponds to the interesting dip in figure 3 at a specific strength of the Coulomb interaction $K_\rho$. It means that by tuning the strengths of the SOIs $\alpha$ and $\beta$ and the Coulomb interaction strength $K_\rho$, the ‘slow’ excitation can be frozen and only the ‘fast’ excitation contributes to electron transport. The slower boson excitation is mainly formed by spin excitations. The static spin susceptibility diverges at the critical value [22].

The orientation-dependent real part of the ac conductance can also be clearly seen from figure 4 where the quantum wires are oriented along two different crystallographic directions $\theta = \pi/4$ and $\theta = 3\pi/4$, respectively. The sharp peaks and dips in the oscillation correspond to the resonance and antiresonance of the symmetric and antisymmetric modes formed in the central wire. From equations (3) and (7), one can see that when $\delta v \neq 0$ the ac conductance oscillation consists of fast and slow oscillations arising from the contributions of two collective excitations propagating along the wire with different velocities $u_1$ and $u_2$, respectively. These two frequencies are determined by the ratio of the velocity $u_{1,2}$ to the length of the wire $L$. Expressing the frequency of the ac voltage $\omega$ in units of $2\pi v_F/L$ as shown in figure 4, we obtain
Figure 4. The ac conductance (in units of $e^2/h$) of a quantum wire with FL leads as a function of $\omega L/(2\pi v_F)$ with fixed $K_\rho = 0.9$. (a) $\alpha = 0.3$, $\beta = 0.2$, $\theta = \pi/4$; (b) $\alpha = 0.3$, $\beta = 0.2$, $\theta = 3\pi/4$; (c) for $\alpha = \beta = 0.2$, $\theta = \pi/4$; (d) $\alpha = \beta = 0.2$, $\theta = 3\pi/4$.

the dimensionless frequencies of the fast and slow ac conductance oscillations $f_{1,2} = v_F/u_{1,2}$, from which the parameter $\delta v$ can be extracted using the relation $\delta v = \sqrt{(u_1^2 + u_2^2)/2 - v_F^2}$. Denoting by $f_{1/2}^\pm$ the two different frequencies at $\theta = \pi/4$ (+), $3\pi/4$ (−), respectively, the ratio of the strength of the RSOI to that of the DSOI is estimated by

$$\frac{\alpha}{\beta} = \frac{\delta v^+ + \delta v^-}{|\delta v^+ - \delta v^-|},$$

with

$$\delta v^\pm = \frac{v_F}{\sqrt{2}} \sqrt{\frac{1}{(f_{1/2}^\pm)^2} + \frac{1}{(f_{1/2}^\mp)^2} - 2.} \tag{10}$$

Using the fast Fourier transformation technique, one can obtain the frequencies and finally obtain the relative strengths of the RSOI $\alpha$ and DSOI $\beta$. Here our analysis shows theoretically that the ac conductance offers an efficient way of measuring the relative strengths of the RSOI and DSOI. It is interesting to note that when $\delta v = 0$, i.e. $\alpha = \beta$ and $\theta = 3\pi/4$, the real part of the ac conductance $\text{Re} G(\omega)$ oscillates as a function of $\omega$ with two different frequencies, $v_F/v_\rho$ and $v_F/v_\sigma$, which are independent of the strengths of SOIs $\alpha$ and $\beta$ (see figure 4(d)). This is because the charge and spin collective excitations are decoupled in this specific case.

Finally, we comment on the experimental feasibility of our proposal. (i) The length of the narrow and clean InAs quantum wire $L = 6 \mu m$ is less than the coherence length $L = 10 \mu m$, which ensures that the ballistic limit can be experimentally achieved [23]. (ii) The state-of-the-art technology makes it possible to fabricate a quantum wire with width $d = 10 \text{ nm}$, so that electrons occupy only the lowest subband. In this case the Coulomb interaction is important and its strength $K_\rho$ ranges from 0.4 to 0.7 in semiconductor quantum wires [24, 25]. (iii) The strengths of the RSOI $\alpha$ and DSOI $\beta$ considered in our calculation

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(in units of $\hbar v_F = 6.58 \times 10^{-11}$ eV m) are approximately $(0.64-4) \times 10^{-11}$ eV m, which are easily achieved in InAs quantum wells. (iv) In figure 4, the frequency of the ac voltage $\omega/2\pi$ (in units of $v_F/L = 1.67$ GHz) ranges from MHz to 12 GHz, which is within the reach of present-day experimental techniques [26, 27]. A very recent experiment measured the ac conductance in Q1D systems [28].

4. Conclusions

In summary, we propose a new scheme for measuring the relative values of the Rashba and Dresselhaus SOIs in interacting quantum wires utilizing the SOI-induced orientation-dependent ac conductance. The orientation dependence of the dc conductance is smeared out in quantum wires with finite length; the ac conductance exhibits interesting oscillations when we change the crystallographic orientation, the strength of the Coulomb interaction and the frequency of the external ac voltage. From the ac conductance one can determine the relative strength of the RSOI and DSOI, which is crucial for understanding spin decoherence and constructing all-electrical spintronic devices.

Acknowledgment

This work was supported by the NSFC under grant numbers 11004017, 60525405, 10874175 and 10947134 and the City University of Hong Kong Strategic Research Grant (project no. 7008108).

References

[1] Žutić I, Fabian J and Sarma S D 2004 Rev. Mod. Phys. 76 323
[2] Winkler R 2003 Spin–Orbit Coupling Effects in Two-Dimensional Electron and Hole Systems (Springer Tracts in Modern Physics) (Berlin: Springer)
[3] Rashba E I 1960 Sov. Phys. Solid State 2 1109
Rashba E I and Sherman E Ya 1988 Phys. Lett. A 129 175
[4] Dresselhaus G 1955 Phys. Rev. 100 580
[5] Yang W and Chang K 2006 Phys. Rev. B 73 113303
[6] Datta S and Das B 1990 Appl. Phys. Lett. 56 665
Schliemann J, Egues J C and Loss D 2003 Phys. Rev. Lett. 90 146801
[7] Ganichev S D et al 2004 Phys. Rev. Lett. 92 256601
[8] Koralek J D, Weber C P, Orenstein J, Bernevig B A, Zhang S C, Mack S and Awschalom D D 2009 Nature 458 610
[9] Meier L, Salis G, Shorubalko I, Gini E, Schön S and Ensslin K 2007 Nat. Phys. 3 650
[10] Schüpers Th, Guzenko V A, Pala M G, Züllicke U, Governale M, Knobbe J and Hardtdegen H 2006 Phys. Rev. B 74 081301
Kunihashi Y, Kohda M and Nitta J 2009 Phys. Rev. Lett. 102 226601
Kallaher R L, Heremans J J, Goel N, Chung S J and Santos M B 2010 Phys. Rev. B 81 035335
[11] Scheid M, Kohda M, Kunihashi Y, Richter K and Nitta J 2008 Phys. Rev. Lett. 101 266401
[12] Luttinger J M 1963 J. Math. Phys. 4 1154
[13] Tarucha S, Honda T and Saku T 1995 Solid State Commun. 94 413
Levy E, Tsukernik A, Karpovski M, Palevski A, Dwir B, Pelucchi E, Rudra A, Kapon E and Oreg Y 2006 Phys. Rev. Lett. 97 196802
[14] Auslaender O M, Steinberg H, Yacoby 1 A, Tserkovnyak 1 Y, Halperin B I, Baldwin K W, Pfeiffer L N and West K W 2005 Science 308 88
[15] Moroz A V, Samokhin K V and Barnes C H W 2000 Phys. Rev. B 62 16900
[16] Gritsev V, Japaridze G, Pletyukhov M and Baeriswyl D 2005 Phys. Rev. Lett. 94 137207
[17] Qimiao S 1998 Phys. Rev. Lett. 81 3191
[18] Governale M and Zülicke U 2002 Phys. Rev. B 66 073311
   Schulz A, De Martino A, Ingenhoven P and Egger R 2009 Phys. Rev. B 79 205432
[19] Dolcini F, Trauzettel B, Safi I and Grabert H 2005 Phys. Rev. B 71 165309
[20] Maslov D L and Stone M 1995 Phys. Rev. B 52 5539
[21] Safi I and Schulz H J 1995 Phys. Rev. B 52 17040
[22] Metzner W and Di Castro C 1993 Phys. Rev. B 47 16107
[23] Honda T, Tarucha S, Saku T and Tokura Y 1995 Japan. J. Appl. Phys. 34 L72
[24] Levy E, Tsukernik A, Karpovski M, Palevski A, Dwir B, Pelucchi E, Rudra A, Kapon E and Oreg Y 2006 Phys. Rev. Lett. 97 196802
[25] Steinberg H, Barak G, Yacoby A, Pfeiffer L N, West K W, Halperin B I and Le Hur K 2008 Nat. Phys. 4 116
[26] Gabelli J and Reulet B 2008 Phys. Rev. Lett. 100 026601
[27] Poo Y, Wu R, Fan X and Xiao J Q 2010 Rev. Sci. Instrum. 81 064701
[28] Yu Z and Burke P J 2005 Nano Lett. 5 1403