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Research Article

Spin Accumulation in a Quantum Wire with Rashba Spin-Orbit Coupling

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We investigate theoretically the spin accumulation in a Rashba spin-orbit coupling quantum wire. Using the scattering matrix approach within the effective free-electron approximation, we have demonstrated the three components of spin polarization. It is found that by a few numerical examples, the two peaks for the out-of-plane spin accumulation $\langle S_z \rangle$ shift to the edges of quantum wire with the increase of propagation modes. The period of intrinsic oscillations $\langle S_{x/y} \rangle$ inversely proportions to the Rashba SOC strength. This effect may be used to differentiate the intrinsic spin accumulation from the extrinsic one.

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Spintronics is a rapidly developing field of research in recent years because spin-based devices are believed to be the new generation of electronic devices [1–3]. Among the research area of spintronics, the spin-orbit coupling (SOC) creates another way to manipulate spins by means of an electric field [3], which has great potential applications in the future. The Rashba spin-orbit coupling (RSOC) effect is found to be very pronounced in some semiconductor layer (e.g., InAs or InGaAs) and its strength can be controlled by gate voltage [4].

The spin polarization in two-dimensional electron gas (2DEG) systems with SOC has been attracting much attention [5–9]. Recently, the experimental observations [10, 11] of spin Hall effect (SHE) in 2DEG with SOC have given evidence of the possibility to control electron spins by an electric field. Wang and Sheng [12] have investigated the current induced local spin polarization due to weak RSOC in a narrow strip. Furthermore, Nikolić et al. [13] have demonstrated that the flow of a longitudinal unpolarized current through a ballistic 2DEG with RSOC will induce a nonequilibrium spin accumulation along the lateral edges. More recently, the out-of-plane spin Hall accumulation effect has been observed in a Rashba bar [14]. Using the nonequilibrium Green's function approach, Wang et al. [15] have studied the intrinsic oscillation of spin accumulation induced by RSOC in a quantum wire (QW) system. However, the work in [12, 14, 15] has only considered the hard-wall potential confined Rashba QWs.

In this paper, using the scattering matrix method within the effective free-electron approximation, we study theoretically the spin accumulation in a parabolic potential confined Rashba QW connected to two normal leads (without SOC) nonadiabatically, and the results are discussed and compared with the previous works [12–15]. The system under consideration is a wide-narrow-wide (WNW) geometry of structure, and the number of propagation modes in the narrow QW and wide leads is controlled by the different strength of parabolic confining potentials. The transverse direction corresponds to x-axis and the longitudinal direction to y-axis. Therefore, the effective mass single-electron Hamiltonian in the QW is

$$H = H_0 + H_{so}, \tag{1}$$

where

$$H_0 = -\frac{\hbar^2}{2m^*}\left(\frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2}\right) + V(x) \tag{2}$$

is the uncoupled QW Hamiltonian while the RSOC Hamiltonian [14, 15] reads

$$H_{so} = i\alpha_R\left(\sigma_y\frac{\partial}{\partial x} - \sigma_x\frac{\partial}{\partial y}\right). \tag{3}$$

In the above Hamiltonians, $V(x) = m^*\omega^2x^2/2$ is the lateral confining potential with m^*, and ω the effective electron mass and oscillator frequency, while α_R denotes the Rashba strength, and $\sigma_{x/y}$ the x/y component of pauli matrix.

Since the Hamiltonian in the leads has no RSOC term, we assume that an unpolarized electron wave is injected from the left lead into the QW. Due to the translation invariance in the longitudinal y-direction, the electron wave function in the leads can be decomposed into two components of a plane wave and a transversally confined wave function for the Nth mode

$$\phi_{N\uparrow\downarrow}(\alpha x) = \frac{e^{-(\alpha x)^2/2}}{\sqrt{2^N N!\sqrt{\pi}}} H_N(\alpha x) \quad (N = 1, 2, \ldots), \quad (4)$$

where $\uparrow(\downarrow)$ represents spin-up (-down) state, $\alpha = (m^*\omega/\hbar)^{1/2}$, and $H_N(\alpha x)$ is the Hermite polynomial. Since no RSOC in the leads, the total electron energy for spin-up or -down electron is same, $E = \hbar^2K_N^2/2m^* + \varepsilon_N$, where the lateral sublevel $\varepsilon_N = (N + 1/2)\hbar\omega$ and the longitudinal wavevector $K_N = \sqrt{2m^*(E - \varepsilon_N)}/\hbar$.

The wave function for the weak Rashba QW is a two-components spinor $\psi(x, y) = \{\psi_\uparrow(x, y)\psi_\downarrow(x, y)\}$ and $\psi_{\uparrow\downarrow}(x, y)$ can be also decomposed into two components for spin-up ($\uparrow$) and spin-down ($\downarrow$) electron, respectively. The corresponding transverse (x direction) wave functions in the perturbation regime [16] can be solved as

$$\varphi_{n\uparrow\downarrow}(\alpha x) = \frac{e^{-(\alpha x)^2/2}}{\sqrt{2^n n!\sqrt{\pi}}} \left\{ H_n(\alpha x) \pm \frac{l_\omega}{l_{\alpha_R}} \left[nH_{n-1}(\alpha x) + \frac{1}{2} H_{n+1}(\alpha x) \right] \right\}, \quad (5)$$

where $l_{\alpha_R} = \hbar^2/(2m^*\alpha_R)$ and $l_\omega = 1/\alpha$ the character length related to the oscillator frequency, and the sign "+/−" corresponding to "$\uparrow$ / $\downarrow$" state. The electron energy in the QW is $E_{\uparrow\downarrow} = \hbar^2k_{n\uparrow\downarrow}^2/2m^* + \varepsilon_{n\uparrow\downarrow}$, where the lateral sublevel of the nth mode is $\varepsilon_{n\uparrow} = (n + 1/2)\hbar\omega + \alpha_R k_{n\uparrow}$ and $\varepsilon_{n\downarrow} = (n + 1/2)\hbar\omega - \alpha_R k_{n\downarrow}$, respectively. Therefore, the longitudinal wavevectors can be expressed as

$$\begin{aligned} k_{n\uparrow}^{\pm}(E_\uparrow) &= \frac{m^*}{\hbar^2} \left[-\alpha_R \pm \sqrt{\alpha_R^2 + 2\hbar^2 \frac{E_\uparrow - (n + 1/2)\hbar\omega}{m^*}} \right], \\ k_{n\downarrow}^{\pm}(E_\downarrow) &= \frac{m^*}{\hbar^2} \left[\alpha_R \pm \sqrt{\alpha_R^2 + 2\hbar^2 \frac{E_\downarrow - (n + 1/2)\hbar\omega}{m^*}} \right] \end{aligned} \quad (6)$$

for spin-up and -down electron, respectively.

In order to calculate the scattering matrix separately for spin-up and -down electron, one needs to match the wave functions at the two interfaces of wire-lead connection. For a symmetric system, one can only calculate the scattering matrix for one interface. The total scattering matrix for the system can be obtained by the symmetry property of the system [17]. Here, we first consider the left interface at $y = 0$ as if without the right interface ($y = L$) for electron incident from left to right. Because of the degeneracy, in the left lead ($y < 0$), the wave functions for spin-up and -down electron can be written as

$$\psi_{\uparrow\downarrow}(x, y) = \phi_{N\uparrow\downarrow}(\alpha x)e^{iK_N y} + \sum_{N'} r_{\uparrow\downarrow N'N}\phi_{N'\uparrow\downarrow}(\alpha x)e^{-iK_{N'} y}, \quad (7)$$

where $r_{\uparrow\downarrow N'N}$ are the coefficients of incident mode N reflecting to mode N'. In the QW ($y > 0$), the wave functions for spin-up and -down electron are nondegenerate

$$\psi_{\uparrow\downarrow}(x, y) = \sum_n t'_{\uparrow\downarrow nN}\varphi_{n\uparrow\downarrow}(\alpha x)e^{ik_{n\uparrow\downarrow}^+ y}, \quad (8)$$

where $\varphi_{n\uparrow\downarrow}$ are given in (5) and $t'_{\uparrow\downarrow nN}$ the coefficients of incident mode N, transmitting to mode n. The continuity of the wave functions at left interface $y = 0$ gives

$$\phi_{N\uparrow\downarrow}(\alpha x) + \sum_{N'} r_{\uparrow\downarrow N'N}\phi_{N'\uparrow\downarrow}(\alpha x) = \sum_n t'_{\uparrow\downarrow nN}\varphi_{n\uparrow\downarrow}(\alpha x). \quad (9)$$

The continuity of the derivative of the wave functions may not hold due to the existence of RSOC [16]. However, the current conservation for the system still holds [18, 19]:

$$\begin{aligned} &\begin{bmatrix} -i\dfrac{\hbar}{m^*}\dfrac{\partial}{\partial y} & 0 \\ 0 & -i\dfrac{\hbar}{m^*}\dfrac{\partial}{\partial y} \end{bmatrix} \begin{bmatrix} \psi_\uparrow(x, y < 0) \\ \psi_\downarrow(x, y < 0) \end{bmatrix} \\ &\quad = \begin{bmatrix} -i\dfrac{\hbar}{m^*}\dfrac{\partial}{\partial y} & \dfrac{\alpha_R}{\hbar} \\ \dfrac{\alpha_R}{\hbar} & -i\dfrac{\hbar}{m^*}\dfrac{\partial}{\partial y} \end{bmatrix} \begin{bmatrix} \psi_\uparrow(x, y \geq 0) \\ \psi_\downarrow(x, y \geq 0) \end{bmatrix}, \end{aligned} \quad (10)$$

hence we can obtain

$$\begin{aligned} &K_N\phi_{N\uparrow\downarrow}(\alpha x) - \sum_{N'} r_{\uparrow\downarrow N'N}K_{N'}\phi_{N'\uparrow\downarrow}(\alpha x) \\ &\quad = \sum_n t'_{\uparrow\downarrow nN}\left[k_{n\uparrow\downarrow}^+\varphi_{n\uparrow\downarrow}(\alpha x) + \varphi_{n\downarrow\uparrow}(\alpha x)/2W\right], \end{aligned} \quad (11)$$

where $W = l_{\alpha_R} = \hbar^2/(2m^*\alpha_R)$ is the effective width of the QW. Multiplying $\phi_{N'\uparrow\downarrow}$ and $\varphi_{n'\uparrow\downarrow}$ with (9) and (11), respectively, and integrating over the range of $-W/2 < x < W$, one obtains the matrices of reflection $r_{\uparrow\downarrow}$ and transmission $t'_{\uparrow\downarrow}$ at the $y = 0$ interface as

$$r_{\uparrow\downarrow} = \lambda_{\uparrow\downarrow}t'_{\uparrow\downarrow} - 1, \quad (12)$$

with

$$\begin{aligned} t'_\uparrow &= \left[\chi_\uparrow^{-1}(\lambda_\uparrow^T K\lambda_\uparrow + p_\uparrow k_\uparrow^+) - (\lambda_\downarrow^T K\lambda_\downarrow + p_\downarrow k_\downarrow^+)^{-1}\chi_\downarrow\right]^{-1} \\ &\quad \times \left[2\chi_\uparrow^{-1}\lambda_\uparrow^T K - 2(\lambda_\downarrow^T K\lambda_\downarrow + p_\downarrow k_\downarrow^+)^{-1}\lambda_\downarrow^T K\right], \\ t'_\downarrow &= \left[\chi_\downarrow^{-1}(\lambda_\downarrow^T K\lambda_\downarrow + p_\downarrow k_\downarrow^+) - (\lambda_\uparrow^T K\lambda_\uparrow + p_\uparrow k_\uparrow^+)^{-1}\chi_\uparrow\right]^{-1} \\ &\quad \times \left[2\chi_\downarrow^{-1}\lambda_\downarrow^T K - 2(\lambda_\uparrow^T K\lambda_\uparrow + p_\uparrow k_\uparrow^+)^{-1}\lambda_\uparrow^T K\right], \end{aligned} \quad (13)$$

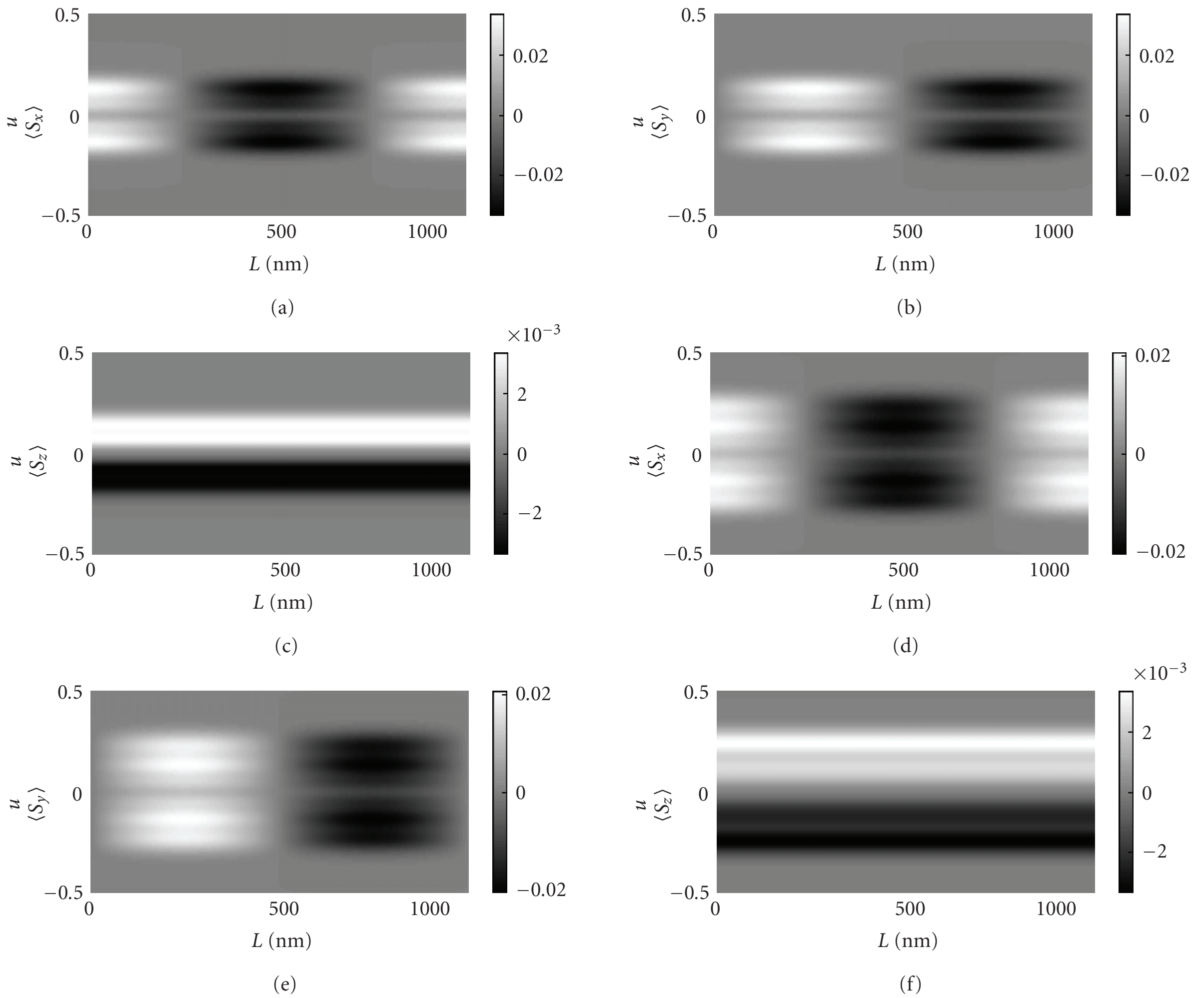

FIGURE 1: The 2D contour of spin polarization components $\langle S_i \rangle$ $(i = x, y, z)$ as functions of transverse position u and length L for $n = 3$ ((a), (b), (c)) and $n = 5$ ((d), (e), (f)) propagation modes QW with $\alpha_R = 0.5 \times 10^{-11}$ eVm.

where the corresponding matrices elements are as follows $p_{nn'\uparrow(\downarrow)} = \int_{-W/2}^{W/2} \varphi_{n\uparrow(\downarrow)}(\alpha x)\varphi_{n'\uparrow(\downarrow)}(\alpha x)dx$, $\chi_{nn'\uparrow(\downarrow)} = (m^* \alpha_R)/\hbar^2 \times \int_{-W/2}^{W/2} \varphi_{n\uparrow(\downarrow)}(\alpha x)\varphi_{n'\downarrow(\uparrow)}(\alpha x)dx$, $\lambda_{Nn\uparrow(\downarrow)} = \int_{-W/2}^{W/2} \phi_{N\uparrow(\downarrow)}(\alpha x) \times \varphi_{n\uparrow(\downarrow)}(\alpha x)dx$, $k_{nn'\uparrow(\downarrow)} = k_{n\uparrow(\downarrow)}\delta_{nn'}$, and $K_{NN'} = K_N \delta_{NN'}$, respectively. Furthermore, an analogous calculation for an incident wave from the right side to the left one at $y = 0$ interfacegives matrices $r'_{\uparrow\downarrow}$ and $t_{\uparrow\downarrow}$. Thus, the scattering matrix $S^{(1)}_{\uparrow\downarrow}$ at $y = 0$ interface can be expressed as

$$S^{(1)}_{\uparrow\downarrow} = \begin{bmatrix} \tilde{r}_{NN'\uparrow\downarrow} & \tilde{t}_{Nn\uparrow\downarrow} \\ \tilde{t}'_{nN\uparrow\downarrow} & \tilde{r}'_{nn'\uparrow\downarrow} \end{bmatrix}, \tag{14}$$

where the normalized wave amplitudes $\tilde{t}'_{nN\uparrow\downarrow} = (k^+_{n\uparrow\downarrow}/K_N)^{1/2} \times t'_{nN\uparrow\downarrow}$, $\tilde{t}_{Nn\uparrow\downarrow} = (K_N/k^-_{n\uparrow\downarrow})^{1/2} t_{Nn\uparrow\downarrow}$, $\tilde{r}_{NN'\uparrow\downarrow} = (K_N/K_{N'})^{1/2} r_{NN'\uparrow\downarrow}$, and $\tilde{r}'_{nn'\uparrow\downarrow} = (k^-_{n\uparrow\downarrow}/k^+_{n\uparrow\downarrow})^{1/2} r'_{nn'\uparrow\downarrow}$. Due to the symmetry of the system, [17] the scattering matrix $S^{(2)}_{\uparrow\downarrow}$ at $y = L$ interface is

$$S^{(2)}_{\uparrow\downarrow} = \begin{bmatrix} \tilde{r}'_{nn'\uparrow\downarrow} & \tilde{t}'_{nN\uparrow\downarrow} \\ \tilde{t}_{Nn\uparrow\downarrow} & \tilde{r}_{NN'\uparrow\downarrow} \end{bmatrix}. \tag{15}$$

Therefore, we can obtain the total wave functions in the QW as

$$\Psi(x, y) = \begin{pmatrix} \psi_\uparrow(x, y) \\ \psi_\downarrow(x, y) \end{pmatrix} = \sum_n \begin{pmatrix} \sum_N \tilde{t}'_{nN\uparrow} e^{ik^+_{n\uparrow} y} \varphi_{n\uparrow}(x) + \sum_{n'} \tilde{r}'_{nn'\uparrow} e^{-ik^+_{n\uparrow}(L-y)} \varphi_{n\uparrow}(x) \\ \sum_N \tilde{t}'_{nN\downarrow} e^{ik^+_{n\downarrow} y} \varphi_{n\downarrow}(x) + \sum_{n'} \tilde{r}'_{nn'\downarrow} e^{-ik^+_{n\downarrow}(L-y)} \varphi_{n\downarrow}(x) \end{pmatrix}, \tag{16}$$

and the local spin polarization [12, 14]

$$\langle S_i(x, y) \rangle = \frac{\hbar}{2} \Psi^\dagger(x, y) \sigma_i \Psi(x, y). \tag{17}$$

In the following, we present some numerical examples of the calculated spin polarization (in units of $\hbar = 1$) in the QW at low temperatures. It should be pointed out that we

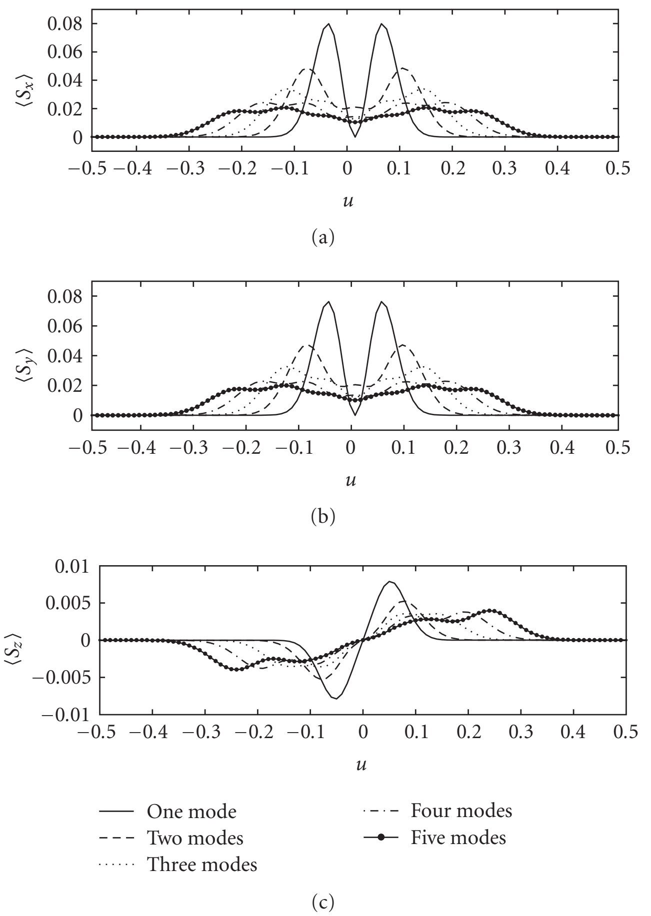

FIGURE 2: The spin polarization components $\langle S_i \rangle$ ($i = x, y, z$) as a function of the transverse position u. The Rashba strength $\alpha_R = 0.5 \times 10^{-11}$ eVm and $\langle S_z \rangle$ indicates the out-of-plane spin accumulation.

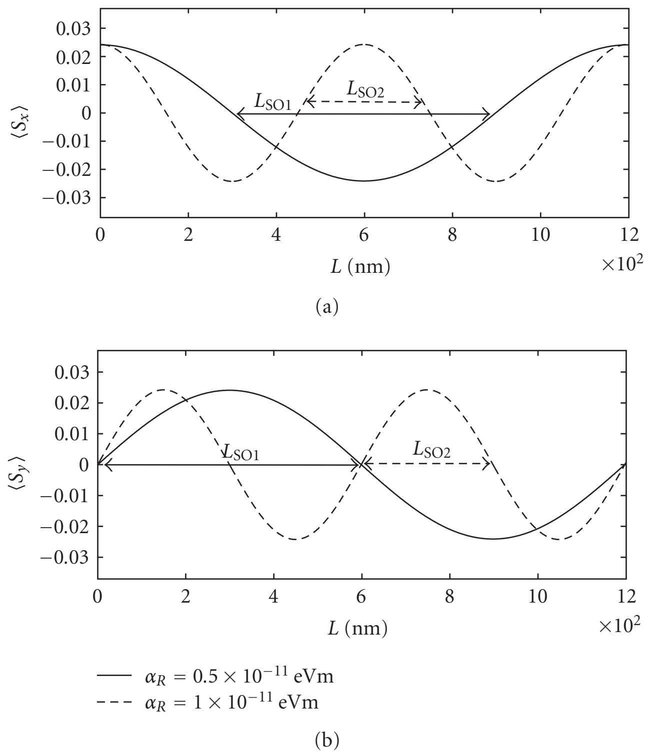

FIGURE 3: The intrinsic oscillations of spin polarization components $\langle S_x \rangle$ and $\langle S_y \rangle$ for $\alpha_R = 0.5 \times 10^{-11}$ eVm (dashed line) and $\alpha_R = 1.0 \times 10^{-11}$ eVm (solid line) with $n = 4$ and the transverse position $u = 0.15$.

only consider the weak RSOC case [16–19], that is, $l_{\omega_0}/l_{\alpha_R} = 2m^*\alpha_R/\hbar^2\alpha_0 = 0.1$ with $l_{\omega_0} = 1/\alpha_0 = (\hbar/m^*\omega_0)^{1/2}$, and the width of two leads supporting number of propagating modes up to $N = 6$. For the case of $\alpha_R = 0.5 \times 10^{-11}$ eVm [7], the incident electron energy is taken as $E = E_{\uparrow\downarrow} = 7\hbar\omega_0 = 0.0367$ eV (for $\alpha_R = 1.0 \times 10^{-11}$ eVm, $E = 0.147$ eV). Our system is a WNW QW ($n < N$); however, the actual width of QW is chosen as $W = l_{\alpha_R}$ and the parabolic confining potential controls the number of propagation modes n in the QW with the corresponding character oscillator frequency length $l_\omega < W$. The electron effective mass is taken as that for InGaAs quantum well $m^* = 0.037m_e$ [8] and the length of QW is chosen as $L = 1200$ nm.

Figure 1 shows the 2D contour of spin polarization components $\langle S_i \rangle$ ($i = x, y, z$) for the cases of three and five propagation modes in the QW. It is noted that the components $\langle S_x \rangle$ and $\langle S_y \rangle$ present periodic oscillations along the longitudinal direction which comes from Rashba spin precession [15] while no oscillations presented for $\langle S_z \rangle$. This is the main difference from the results of [14] due to the different confining potential.

The polarization components $\langle S_i \rangle$ versus the transverse position $u = x/W$ for different propagation modes in the QW are demonstrated in Figure 2, with the longitudinal positions for $\langle S_x \rangle$ and $\langle S_{y/z} \rangle$ as 0 and 300 nm, respectively. More peaks and extended polarization width toward two edges of the QW have been shown with the increasing modes n. The component $\langle S_z \rangle$ presents two inverse-direction peaks (shifting to the two edges of the QW), consequently, produces the out-of-plane spin accumulation in the QW. The spin accumulation is formed with the spin distribution polarized along the z direction due to the interaction between the pseudomagnetic field arising from RSOC and the accelerating electrons [20, 21]. Moreover, the net spin polarization rather than in-plane spin accumulation is illustrated for $\langle S_{x/y} \rangle$ while the present results of $\langle S_z \rangle$ are qualitatively consistent with [12, 14].

Furthermore, in Figure 3 we plot the longitudinal profile of spin polarization components $\langle S_x \rangle$ and $\langle S_y \rangle$ at a fixed transverse position $u = 0.15$. As can be seen, due to the Rashba spin precession along the longitudinal position of QW, the period of the oscillations (with $\pi/2$ phase separation) for Rashba strength $\alpha_R = 1.0 \times 10^{-11}$ eVm is half that of the 0.5×10^{-11} eVm case. One notes that the Rashba precession lengths L_{SO1} and L_{SO2} are 600 nm and 300 nm, respectively, which are consistent with the results ($L_{SO} = \pi\hbar^2/(2m^*\alpha_R)$) in [14, 22]. It should be pointed out that the presented results may be observable experimentally to differentiate the intrinsic and extrinsic spin accumulation.

In conclusion, we have studied the spin polarization in a Rashba QW confined by a parabolic potential. It is found that the two inverse direction peaks of the out-of-plane spin accumulation $\langle S_z \rangle$ shift to the two edges of the QW with the increasing propagation modes. Furthermore, the intrinsic oscillation of spin polarization due to Rashba spin precession can be used to identify the intrinsic spin accumulation by changing the RSOC strength.

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Magnetoelectricity: A Bridge across the Disciplines of Nanotechnology, Materials Science, and Electromagnetism

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The idea of magnetoelectricity has been around for decades, and several materials and concepts have been studied; however, recent discoveries of compounds in which the magnetoelectric (ME) coupling is strong have sparked a renaissance. ME symmetry allows magnetization (resp., polarization) to be induced by the application of an electric (resp., magnetic) field. The initial impetus driving ME materials was the prospect of employing a single field excitation, for example, voltage gradient, to tune both polarization and magnetization characteristics. Significant progress in ME nanostructures has been made in recent years, mainly driven by materialist use in new class of multifunctional device applications, for example, electric field-controlled magnetic data storage, antenna miniaturization using passive substrates, and field control for spintronics applications. Manipulation of electric polarizations by external magnetic fields (and vice versa) has been demonstrated in some of single-phase multiferroic materials, for example, $BiFeO_3$. In this case, the intrinsic ME effect arises from the coupling between electric and magnetic order parameters. Intrinsic ME coupling occurs in compounds in which time-reversal and space-inversion symmetries are absent. A significant ME effect is also observed in the case of two-phase (thin films) composite multiferroics made by combining ferroelectric and ferromagnetic substances together, for example, $BaTiO_3/CoFe_2O_4$, and lead-zirconate-titanate (PZT)/ferrite (or Terfenol-D). In these previous studies, it is commonly believed that elastic coupling amongst the ferroelectric and magnetic order parameters is at the origin of the ME effect, that is, an electric field E induces strain in the ferroelectric; this strain is passed on to the ferromagnet, where it causes a net magnetization M in the sample. A basic understanding of the distributions of electric, magnetic, and internal stress fields within the material allows one to manipulate the relative stability of domain structures. One would expect that the ME effect is large if the coupling at the interface is large; thus, composites with large surface area (such as multilayer systems and nanogranular mixtures) and strongly ferroelastic components are especially effective. The theoretical understanding of the microscopic mechanism that causes the strong coupling between the electric and magnetic degrees of freedom is quite limited. We therefore anticipate that ME nanostructures may open up a whole new dimension for a wide range of materials having cost effective, small size planar, and broadband multifunctional device applications.

Papers are solicited in, but not limited to, the following areas:

- Physical mechanisms that control the ME property in layered and granular nanostructures
- Multiferroic composites
- Multifunctional applications of nanostructures
- Giant ME effects
- Magnetoelectric interactions at electromechanical resonance
- Magnetization and polarization tuning by external fields
- Coexistence of magnetic and ferroelectric orders
- Static and dynamic magnetoelasticity
- Crystallographic engineering by assembling ME building blocks
- ME effect in superconductors

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Manuscript Due	December 1, 2009
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Advances in Condensed Matter Physics

Special Issue on Phonons and Electron Correlations in High-Temperature and Other Novel Superconductors

Call for Papers

It has been now over 20 years since the discovery of the first high-temperature superconductor by Georg Bednorz and Alex Müller and yet, despite intensive effort, no universally accepted theory exists about the origin of superconductivity in cuprates. The absence of consensus on the physics of cuprate superconductors and the recent discovery of iron-based compounds with high transition temperatures have re-emphasized the fundamental importance of understanding the origin of high-temperature superconductivity. First-principles calculations based on density functional theory (DFT) often predict a rather weak electron-phonon interaction (EPI) insufficient to explain high transition temperatures. A number of researchers are of the opinion that the EPI may be considerably enhanced by correlation effects beyond DFT. Others maintain that the repulsive electron-electron interaction in novel superconductors is pairing and provides high transition temperatures without phonons. On the other hand, some recent studies using numerical techniques cast doubt that simple repulsive models can account for high-temperature superconductivity. Therefore, it seems plausible that the true origin of high-temperature superconductivity could be found in a proper combination of strong electron-electron correlations with a significant EPI.

This Special Issue will become an international forum for researchers to summarize recent developments in the field, with a special emphasis on the results in high-temperature superconductors and some other related materials that combine sizeable electron-phonon coupling with strong correlations. We invite authors to present original research papers as well as summarizing overviews stimulating the continuing efforts to understand high-temperature and other unconventional superconductors. Potential topics include, but are not limited to:

- Phonon spectroscopes of cuprates and related compounds
- Experimental evidence for electron-phonon interactions
- First-principles calculations of EPIs
- Strong-coupling extensions of the BCS-Eliashberg theory including polarons and bipolarons
- Interplay of electron-electron and electron-phonon interactions (Holstein-Hubbard, Fröhlich-Hubbard, tJ-Holstein, and similar models)
- Phase separation of correlated electrons
- EPI effects in fullerenes, MgB_2, ruthenates, ferropnictides, and other novel noncuprate superconductors
- Routes to higher temperature superconductivity

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