Suspended Nanowires as Mechanically Controlled Rashba Spin Splitters

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Suspended nanowires are shown to provide mechanically controlled coherent mixing or splitting of the spin states of transmitted electrons, caused by the Rashba spin-orbit interaction. The sensitivity of the latter to mechanical bending makes the wire a tunable nanoelectromechanical weak link between reservoirs. When the reservoirs are populated with misbalanced ''spin-up and spin-down'' electrons, the wire becomes a source of split spin currents, which are not associated with electric charge transfer and which do not depend on temperature or driving voltages. The mechanical vibrations of the bended wires allow for additional tunability of these splitters by applying a magnetic field and varying the temperature. Clean metallic carbon nanotubes of a few microns length are good candidates for generating spin conductance of the same order as the charge conductance (divided by e^2) which would have been induced by electric driving voltages.

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Introduction.—The lack of screening and the wavy nature of the electrons together with the ensuing interference effects determine a large variety of Coulomb-correlation and quantum-coherence phenomena in quantum wires and dots. The electronic spin, being weakly coupled to other degrees of freedom in bulk materials, becomes an ''active player'' due to the enhanced spin-orbit interaction induced by the Rashba effect [1] in these low-dimensional structures [2,3]. This interaction can be also modified experimentally [4–6]. The quantum-coherence control of spin-related devices and the spatial transfer of the electron spins are among the most challenging tasks of current spintronics, as they can bring up new functionalities. Thus, e.g., quantum interference of electronic waves in multiply connected devices was predicted to be sensitive to the electronic spin, leading to spin filtering in electronic transport [7].

In charge transport, electronic beam splitters (e.g., by tunnel barriers) are key ingredients in interference-based devices. In this Letter we propose that tunnel-barrier scatterers may serve as coherent splitters of the electronic spin when the tunneling electrons also undergo spin (Rashba) scattering. This allows us to map various interference based phenomena in charge transport onto electronic spin transportation. Such spin splitters can be readily made functional by adding to them a mechanical degree of freedom, which serves to control their geometrical configuration in space, to which the Rashba interaction is quite sensitive. Because of this, one achieves mechanical coherent control and mechanical tuning of the spin filters [8].

We suggest that a suspended nanowire, acting as a weak link between two electronic reservoirs, is a good candidate for such a Rashba spin splitter (see Fig. 1). The amount of spin splitting, brought about by the Rashba interaction on the wire, is determined by the spin-orbit coupling and as such can be controlled by bending the wire. This can be mechanically tuned, by exploiting a break junction as a substrate for the wire (see Fig. 1) or by electrically inducing a Coulomb interaction between the wire and an STM tip electrode (also displayed in Fig. 1). This Rashba scatterer is localized on the nanowire, and serves as a pointlike scatterer in momentum-spin space for the electrons incident from the bulky leads. When there is a spin imbalance population in one of the leads (or both), and the Rashba spin splitter is activated (i.e., the weak link is open for electronic propagation) spin currents are generated and are injected from the pointlike scatterer to the leads. Thus the

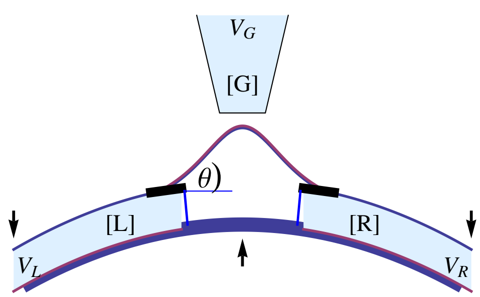

FIG. 1 (color online). A break junction supporting a nanowire of length d (possibly a carbon nanotube), attached by tunnel contacts to two biased electrodes ([L] and [R]). The small vibrations of the wire induce oscillations in the angle θ around some value θ_0. The upper electrode ([G]) is an STM tip biased differently. The Rashba interaction can be controlled via the bending angle θ of the wire. The latter can be modified both mechanically, by loads (shown by the arrows) applied to the substrate and electrically, by biasing the STM.

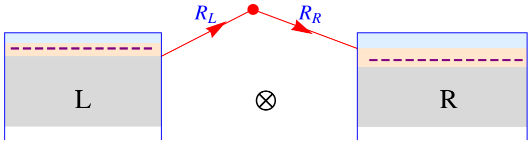

FIG. 2 (color online). Schematic geometry used for calculating the spin-orbit coupling dependence of the tunneling amplitude. A localized level is tunnel coupled to left (L) and right (R) electronic electrodes with possibly different chemical potentials $\mu_{L\sigma}$ and $\mu_{R\sigma}$. The setup lies in the x-y plane; a magnetic field applied along $\hat{\mathbf{z}}$ is shown by ⊗. The setup corresponds to a configuration in which the wire is controlled only mechanically, and the STM is not shown.

Rashba splitter redistributes the spin populations between the leads. This source of the spin currents need not be accompanied by transfer of electronic charges [9].

Such a coherent scatterer, whose scattering matrix can be "designed" at will by tuning controllably the geometry, can be realized in electric weak links based on clean carbon nanotubes (CNT). Carbon nanotubes have a significant Rashba spin-orbit coupling [2,4,6]. Moreover, CNT's are known to have quite long mean-free paths (longer for *suspended* tubes that for the non-bended ones), allowing for experimental detection of interference-based phenomena (e.g., Fabry-Perot interference patterns) [10].

Further tunability of the Rashba spin splitter can be achieved by switching on an external magnetic field, coupled to the wire through the Aharonov-Bohm effect [11]. This is accomplished by quantum-coherent displacements of the wire, which generate a temperature dependence in the Aharonov-Bohm magnetic flux (through an effective area) [12]. Generally, a large mechanical deformability of nanostructures, originating from their composite nature complemented by the strong Coulomb forces accompanying single-electron charge transfer, offer an additional functionality of electronic nanodevices [13,14]. Indeed, coherent nanovibrations in suspended nanostructures, with frequency in the gigahertz range, were detected experimentally [15].

The transmission amplitude through a Rashba scatterer.—The model system exploited in the calculations is depicted in Fig. 2. There, the nanowire is replaced by a quantum dot (a widely accepted picture, see Ref. [10]), which has a single level (of energy ϵ_0), and which vibrates in the direction perpendicular to the wire in the junction plane. The leads are modeled by free electron gases and are firmly coupled to left and right reservoirs, of chemical potentials $\mu_{L\sigma}$ and $\mu_{R\sigma}$, respectively, allowing for spin-polarized charge carriers. Here, σ denotes the spin index; the spin-quantization axis (assumed to be the same for both reservoirs) depends on the spin imbalance in the reservoirs and will be specified below. The electronic populations in the reservoirs are thus

$$f_{L(R)\sigma}(\epsilon_{k(p)}) = [e^{\beta(\epsilon_{k(p)}-\mu_{L(R)\sigma})} + 1]^{-1}, \tag{1}$$

with $\beta^{-1} = k_B T$. The electron gas states in the left (right) lead are indexed by k (p) and have energies ϵ_k (ϵ_p). Below we denote by $c_{k\sigma}$ ($c_{p\sigma}$) the annihilation operators for the leads, and by $c_{0\sigma}$ that for the localized level [16].

The linear Rashba interaction manifests itself as a phase factor on the tunneling amplitude [17]. In the geometry of Fig. 2, this phase is induced by an electric field perpendicular to the x-y plane, and is given by $\alpha \mathbf{R} \times \boldsymbol{\sigma} \cdot \hat{\mathbf{z}}$, where α denotes the strength of the spin-orbit interaction (in units of inverse length [18]), and $\boldsymbol{\sigma}$ is the vector of the Pauli matrices. Quite generally, $\mathbf{R}_L = \{x_L, y_L\}$ for the left tunnel coupling and $\mathbf{R}_R = \{x_R, -y_R\}$ for the right one, where both radius vectors $\mathbf{R}_L$ and $\mathbf{R}_R$ are functions of the vibrational degrees of freedom (as specified in the following). The quantum vibrations of the wire which modify the bending angle, make the electronic motion effectively two dimensional. This leads to the possibility of manipulating the junction via the Aharonov-Bohm effect, by applying a magnetic field which imposes a further phase on the tunneling amplitudes $\phi_{L(R)} = -(\pi/\Phi_0)(Hx_{L(R)}y_{L(R)})$, where H is the magnetic field and Φ_0 is the flux quantum (a factor of order unity is absorbed in H [12]).

It follows that the tunneling Hamiltonian between the localized level and the leads takes the form

$$\mathcal{H}_{\rm tun} = \sum_{k,\sigma,\sigma'} (V_{k\sigma\sigma'} c^\dagger_{0\sigma} c_{k\sigma'} + \text{H.c.}) + \sum_{p,\sigma,\sigma'} (V_{p\sigma\sigma'} c^\dagger_{p\sigma} c_{0\sigma'} + \text{H.c.}). \tag{2}$$

The tunneling amplitudes are (operators in spin and vibration spaces)

$$V_{k(p)} = -J_{L(R)} \exp[-i\psi_{L(R)}], \tag{3}$$

where

$$\begin{aligned} \psi_L &= \phi_L - \alpha(x_L\sigma_y - y_L\sigma_x), \\ \psi_R &= \phi_R - \alpha(x_R\sigma_y + y_R\sigma_x). \end{aligned} \tag{4}$$

We consider a nonresonant case, where the localized level is far above the energies of the occupied states in both leads (i.e., no energy level on the wire is close enough to ϵ_0 to be involved in inelastic tunneling via a real state). This allows us to exploit the tunneling as an expansion parameter [12] and to preform a unitary transformation which replaces the wire by an effective direct tunneling between the leads through virtual states

$$\mathcal{H}^e_{\rm tun} = \sum_{k,p} (c^\dagger_k W^\dagger_{kp} c_p + \text{H.c.}), \tag{5}$$

with (using matrix notations in spin space)

$$W^\dagger_{kp} = \frac{1}{2}\left(\frac{1}{\epsilon_p - \epsilon_0} + \frac{1}{\epsilon_k - \epsilon_0}\right) V^\dagger_k V^\dagger_p. \tag{6}$$

A straightforward calculation [19] now yields that the spin-polarized particle flux emerging from the left lead is

$$I_{L\sigma} = \int_0^\infty d\tau \sum_{k,p,\sigma'} \{ f_{R\sigma'}(\epsilon_p)[1 - f_{L\sigma}(\epsilon_k)] \langle e^{i(\epsilon_k - \epsilon_p)\tau} [W_{pk}]_{\sigma'\sigma} [W^\dagger_{kp}(\tau)]_{\sigma\sigma'} + (\tau \to -\tau) \rangle$$
$$- f_{L\sigma}(\epsilon_k)[1 - f_{R\sigma'}(\epsilon_p)] \langle e^{i(\epsilon_p - \epsilon_k)\tau} [W_{kp}]_{\sigma\sigma'} [W^\dagger_{pk}(\tau)]_{\sigma'\sigma} + (\tau \to -\tau) \rangle \}, \quad (7)$$

where $\langle\rangle$ denotes thermal averaging over the vibrations and over the time evolution with respect to the (free) Hamiltonians of the leads and the vibrations. Assuming that the k, p dependence of the amplitudes may be ignored, and adopting the Einstein model for the description of the vibrations in the variable θ (see below), one readily obtains [19]

$$I_{L\sigma} = \frac{\Gamma_L \Gamma_R}{2\pi\epsilon_0^2} \sum_{\sigma'} \sum_{n,n'=0}^{\infty} P(n) |\langle n| [e^{-i\psi_R} e^{-i\psi_L}]_{\sigma'\sigma} |n'\rangle|^2$$
$$\times (1 - e^{\beta(\mu_{L\sigma} - \mu_{R\sigma'})}) \frac{\mu_{L\sigma} - \mu_{R\sigma'} + (n' - n)\omega}{e^{\beta[\mu_{L\sigma} - \mu_{R\sigma'} + (n'-n)\omega]} - 1}, \quad (8)$$

where n is the vibrations' quantum number, $P(n) = (1 - \exp[-\beta\omega]) \exp[-n\beta\omega]$, and ω is the vibrations' frequency ($\Gamma_{L(R)}$ are the usual partial widths induced on ϵ_0 by the coupling to the leads). The particle flux emerging from the right lead is obtained from Eq. (7) upon interchanging the roles of the left and right sides of the junction, with $\sum_\sigma (I_{L\sigma} + I_{R\sigma}) = 0$, as required by charge conservation. One notes [see Eq. (4)] that while the phase due to the magnetic field disappears in the absence of the vibrations, this is not so for the spin-orbit-phase (as ψ_L and ψ_R do not commute).

The Rashba scatterer as a spin source.—Combining the expressions for the incoming spin fluxes [Eq. (8) and the corresponding one for $I_{R\sigma}$] yields a net spin current, which is injected from the Rashba scatterer into the leads. Therefore, the scatterer can be viewed as a source of spin current, which is maintained when the leads have imbalanced populations. This spin current is defined as

$$J_{\rm spin} \equiv \sum_\sigma \sigma J_{{\rm spin},\sigma} = \sum_\sigma \sigma (I_{L\sigma} + I_{R\sigma}), \quad (9)$$

and it tends to diminish the spin imbalance in the leads, through spin-flip transitions induced by the Rashba interaction. In the limit of weak tunneling, we expect the spin imbalance to be kept constant in time by injecting spin-polarized electrons into the reservoirs, so that the (spin-dependent) chemical potentials do not vary.

The explicit expressions for the two spin currents yield dramatic consequences [20]. (i) Independent of the choice of the spin-quantization axis, $J_{{\rm spin},\sigma}$ is given solely by the term with $\sigma' = \bar{\sigma}$ in the spin sums of Eq. (8) and the corresponding one for $I_{R\sigma}$ ($\bar{\sigma}$ is the spin projection opposite to σ), which implies that only the off-diagonal (in spin space) amplitudes contribute. (ii) Adopting the plausible geometry, $y_L = y_R = (d/2)\sin(\theta)$ and $x_L = x_R = (d/2)\cos(\theta)$, where d is the wire length and θ, which vibrates around θ_0, is defined in Fig. 1, one finds that

$$e^{-i\psi_R} e^{-i\psi_L} = e^{i(\pi H d^2 \sin(2\theta)/4\Phi_0)} (1 - 2\cos^2(\theta)\sin^2(\alpha d/2)$$
$$+ i\sigma_y \sin(\alpha d)\cos(\theta)$$
$$- i\sigma_z \sin(2\theta)\sin^2(\alpha d/2)). \quad (10)$$

This result is independent of the choice of the spin polarizations in the leads, and *does not involve* σ_x. (iii) As Eq. (10) indicates, spin flip will be realized in our device for any orientation of the leads' polarization. Furthermore, if the angle θ vibrates about a nonzero average value θ_0 then both terms on the second and third lines in Eq. (10) yield spin flips even for the nonvibrating wire. In this respect, the spin-orbit splitting effect is very different from that of the Aharonov-Bohm field, which requires a finite area and therefore in our setup is entirely caused by the mechanical vibrations. In the special case $\theta_0 = 0$, the second term there does not contribute for the nonvibrating wire, and then one has spin flips only if the polarization is in the x-z plane. To be concrete, below we present explicit results for a quantization axis along $\hat{\mathbf{z}}$.

Let the chemical potentials of the two leads be

$$\mu_{L,R\uparrow} = \mu_{L,R} + \frac{U_{L,R}}{2}, \qquad \mu_{L,R\downarrow} = \mu_{L,R} - \frac{U_{L,R}}{2}, \quad (11)$$

such that the bias is given by $\Delta\mu = \mu_L - \mu_R$ while the amount of polarization in each of the leads is determined by $U_{L(R)}$. Equation (8) then yields

$$-J_{{\rm spin},\downarrow} = J_{{\rm spin},\uparrow} = \sin^2(\alpha d) \frac{\Gamma_L \Gamma_R}{2\pi\epsilon_0^2} \sum_{n,n'} P(n) |\langle n| e^{i(\pi H d^2/4\Phi_0)\sin(2\theta)} \cos(\theta) |n'\rangle|^2$$
$$\times \left(\frac{[1 - e^{\beta(\Delta\mu + U)}][\Delta\mu + U + (n' - n)\omega]}{e^{\beta[\Delta\mu + U + (n'-n)\omega]} - 1} + \frac{[e^{\beta(\Delta\mu - U)} - 1][\Delta\mu - U + (n' - n)\omega]}{e^{\beta[\Delta\mu - U + (n'-n)\omega]} - 1} \right), \quad (12)$$

where $U = (U_L + U_R)/2$. Clearly, the spin intensity vanishes for non-polarized leads, for which $U = 0$. One also observes that when the vibrations are ignored, i.e., when $\theta = \theta_0$ (and, consequently, only the $n = n'$ terms survives) the Aharonov-Bohm flux drops out, and

$$J_{{\rm spin},\sigma} = G_0 U \sin^2(\alpha d) \cos^2(\theta_0), \quad (13)$$

where $G_0 = \Gamma_L \Gamma_R / (\pi\epsilon_0^2)$ is the Landauer zero-field electric conductance (divided by e^2, e being the electronic charge). Finally, when the spin intensity is linear in the chemical potentials (i.e., in the linear-response regime) it loses its dependence on the bias voltage and becomes

$$J_{\mathrm{spin},\uparrow} = -UG_{\mathrm{spin}}, \tag{14}$$

with the "spin conductance"

$$G_{\mathrm{spin}} = G_0 \sin^2(\alpha d) \sum_{n=0}^{\infty} \sum_{\ell=1}^{\infty} P(n) \times |\langle n| e^{i(\pi H d^2/4\Phi_0)\sin(2\theta)} \cos(\theta) |n+\ell\rangle|^2 \frac{2\ell\beta\omega}{e^{\beta\ell\omega}-1}. \tag{15}$$

Our final expression for the amount of spin intensity is obtained upon expanding $\theta = \theta_0 + \Delta\theta = \theta_0 + (a_0\cos(\theta_0)/d)(b+b^\dagger)$, where b ($b^\dagger$) is the destruction (creation) operator of the vibrations, and a_0 is the amplitude of the zero-point oscillations. Equation (15) then becomes

$$G_{\mathrm{spin}} = \sin^2(\alpha d)\cos^2(\theta_0) \times \left[G_0 \sum_{n=0}^{\infty} \sum_{\ell=1}^{\infty} |\langle n| e^{i(H/H_0)(b+b^\dagger)} |n+\ell\rangle|^2 \frac{2P(n)\ell\beta\omega}{e^{\beta\ell\omega}-1} \right], \tag{16}$$

with the magnetic-field scale given by $H_0 = \sqrt{2}\Phi_0/[\pi d a_0 \cos(\theta_0)\cos(2\theta_0)]$. Interestingly enough, the expression in the square brackets of Eq. (16) is exactly the magnetoconductance (divided by e^2) of the wire, as analyzed in Ref. [12]. Hence, we may build on their results to obtain for the spin-intensity admittance the low- and high-temperature limits

$$\frac{G_{\mathrm{spin}}/G_0}{\sin^2(\alpha d)\cos^2(\theta_0)} = \begin{cases} 1 - \frac{\beta\omega}{6}\frac{H^2}{H_0^2} & \beta\omega \ll 1 \\ \exp[-H^2/H_0^2], & \beta\omega \gg 1. \end{cases} \tag{17}$$

Discussion.—In conclusion, we have proposed that electromechanically tunable interference of waves of electronic spins can be achieved in nanostructures with spatially localized spin-orbit interaction. Electric weak links in a mechanically controllable geometry enable one to exploit the Rashba spin-orbit interaction to split electronic spins and to induce spin currents in polarized conductors. These currents are not associated with electric charge transportation. The Rashba spin splitter is characterized by a scattering matrix that can be "designed" at will, by mechanically tuning the nanowire.

Carbon nanotubes are in particular suitable for realizing the Rashba spin splitter. The energy gap induced by the spin-orbit coupling in them is 0.37 meV, making the strength α on the order of 10^4 cm^{-1} [4]. For wire lengths of the order of micrometers, αd is of order unity, and then G_{spin} is of the same order as the Landauer conductance (which determines the response of electric currents to electric driving voltages). It is then possible to tune the spin currents by an external electric field (which controls the Rashba coupling).

Any experimental detection of spin patterns of electrons flowing through a nanotube would be an appropriate method to monitor the spin current injected from the Rashba splitter. Spin-dependent tunneling is one possibility. When the leads are spin polarized, the densities of states at the Fermi energy of the left (right) lead, $\mathcal{N}^{\sigma}_{L(R)}$, depend on the spin direction. Then the electric conductance, G_0, becomes a function of the mechanical angle θ_0 due to the Rashba-induced spin-flip transitions, leading to

$$J_{\mathrm{spin}}(\theta_0) \propto U\left[G_0(\theta_0) - G_0\left(\frac{\pi}{2}\right)\right] \left| \frac{\mathcal{N}^{\sigma}_L \mathcal{N}^{\bar\sigma}_R - \mathcal{N}^{\bar\sigma}_L \mathcal{N}^{\sigma}_R}{\mathcal{N}^{\sigma}_L \mathcal{N}^{\bar\sigma}_R + \mathcal{N}^{\bar\sigma}_L \mathcal{N}^{\sigma}_R} \right|.$$

Thus an electric measurement of $G_0(\theta_0)$ can detect the spin current. Another scheme which avoids a voltage drop is to exploit the Rashba splitter as a superconducting weak link. The full consideration of this setup is beyond the present scope. However, having in mind the well-known π shift in the Josephson current as a function of the superconducting phase difference [21] resulting from spin flips induced by impurities in the tunnel barrier, one may hope to observe mechanically controlled phase shifts in the Josephson current arising from the Rashba splitting.

The possibility to electrically and mechanically activate pointlike sources of spin-polarized currents with controllable orientations of the spin polarization opens a new route to study spin-related interference in more complicated arrays and networks. It couples such phenomena with electronic transport switching caused by e.g., Coulomb blockade or microwave activation.

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