

# Controlling spin in an electronic interferometer with spin-active interfaces

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(Dated: November 29, 2005)

We consider electronic current transport through a ballistic one-dimensional quantum wire connected to two ferromagnetic leads. We study the effects of the *spin-dependence* of interfacial phase shifts (SDIPS) acquired by electrons upon scattering at the boundaries of the wire. The SDIPS produces a spin splitting of the wire resonant energies which is tunable with the gate voltage and the angle between the ferromagnetic polarizations. This property could be used for manipulating spins. In particular, it leads to a giant magnetoresistance effect with a sign tunable with the gate voltage and the magnetic field applied to the wire.

PACS numbers: 73.23.-b, 75.75.+a, 85.75.-d

The quantum mechanical spin degree of freedom is now widely exploited to control current transport in electronic devices [1]. However, one major functionality to be explored is the electric field control of spin. In the context of a future spin electronics or spintronics, this would allow to build the counterpart of the field-effect transistor (FET), namely the spin-FET, in which spin transport would be controlled through an electrostatic gate [2, 3]. In devices where single spins are used to encode quantum information, this property should also allow to perform single quantum bit operations by using effective magnetic fields which would be locally controllable with the gate electrostatic potential [4]. Among the potential candidates for implementing the electric field control of the spin dynamics, spin-orbit coupling seems a natural choice [2]. However, whether it is possible to use spin-orbit coupling to make spin-FETs or spin qubits is still an open question [5].

In this context, it is crucial to take into account that the interface between a ferromagnet and a non-magnetic material can scatter electrons with spin parallel or antiparallel to the magnetization of the ferromagnet with different phase shifts. This *spin-dependence* of interfacial phase shifts (SDIPS) can modify significantly the behavior of hybrid circuits. First, the SDIPS implies that spins non-collinear to the magnetization of the ferromagnet precess during the interfacial scattering, like the polarization of light rotates upon crossing a birefringent medium. This precession is expected to increase the current through diffusive F/normal metal/F spin valves [6] when the magnetizations of the two F electrodes are non-collinear. The same phenomenon is predicted to occur in F/Luttinger liquid/F [7] and F/Coulomb blockade island/F [8] spin valves in the incoherent limit. In collinear configurations, precession effects are not relevant, but the SDIPS can affect mesoscopic coherence phenomena. For instance, in superconducting/F hybrid systems [9–11], the SDIPS manifests itself by introducing a phase shift between electrons and holes coupled coherently by Andreev reflections. References [9] and [11] have identified experimental signatures of this effect

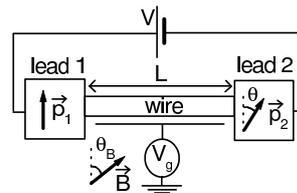


FIG. 1: Electrical diagram of a ballistic wire of length  $L$  connected to ferromagnetic leads 1 and 2 with polarizations  $\vec{p}_1$  and  $\vec{p}_2$ . The wire is capacitively coupled to a gate voltage source  $V_g$ . A magnetic field  $\vec{B}$  is applied to the circuit. We assume that  $\vec{p}_1$ ,  $\vec{p}_2$  and  $\vec{B}$  are coplanar, with angles  $\theta = (\vec{p}_1, \vec{p}_2)$  and  $\theta_B = (\vec{p}_1, \vec{B})$ .

in the data of [12] and [13], respectively. However, the SDIPS has not been shown to affect normal systems in collinear configurations up to now.

The purpose of this Letter is to show that the SDIPS can indeed affect normal systems in collinear configurations, leading to properties which could be used for controlling spins in the context of spintronics and quantum computing. We consider a non-interacting one-dimensional ballistic wire contacted by two ferromagnetic leads. In this system, Fabry-Perot-like energy resonances occur due to size quantization, as observed experimentally for instance in carbon nanotubes contacted by normal electrodes [14]. We show that the SDIPS modifies qualitatively the behavior of this device even in the collinear configuration due to coherent multiple reflections. More precisely, the SDIPS leads to a spin-splitting of the resonant energies which is tunable with the gate voltage and the angle between the ferromagnetic polarizations. This property can be used for manipulating the spin degree of freedom. In particular, the SDIPS-induced spin-splitting can lead to a giant magnetoresistance effect with a sign tunable with the gate voltage and the magnetic field, which should allow to build very efficient spin-FETs.

We consider a single-channel ballistic wire of length  $L$  contacted by ferromagnetic leads 1 and 2 (Fig. 1).

This wire is capacitively coupled to a gate biased at a voltage  $V_g$ , which allows to tune its chemical potential. The directions of the magnetic polarizations  $\vec{p}_1$  and  $\vec{p}_2$  of leads 1 and 2 form an angle  $\theta = (\vec{p}_1, \vec{p}_2)$ . The spin states parallel (antiparallel) to  $\vec{p}_n$  are denoted  $\uparrow_n$  ( $\downarrow_n$ ) in general expressions, or  $u(d)$  in expressions referring explicitly to lead  $n$  only. The wire is subject to a DC magnetic field  $\vec{B}$  coplanar to  $\vec{p}_1$  and  $\vec{p}_2$ , with  $\theta_B = (\vec{p}_1, \vec{B})$ . Following [15], we use a scattering description [16] in which an interface L/R is described by a scattering matrix  $S$  such that  $[\hat{a}_{L,-}^1, \hat{a}_{R,+}^1, \hat{a}_{L,-}^1, \hat{a}_{R,+}^1]^t = S[\hat{a}_{L,+}^1, \hat{a}_{R,-}^1, \hat{a}_{L,+}^1, \hat{a}_{R,-}^1]^t$ , with  $\hat{a}_{L[R],+}^s$  the annihilation operator associated to the right-going (left-going) electronic state with spin  $s$  at the left[right] side of the interface (we use spin space  $\{\uparrow_1, \downarrow_1\}$  for defining  $S$ ). In the low  $V$  limit, the electrostatic potential of the wire is  $\alpha V_g$ , with  $\alpha$  the ratio between the gate capacitance and the total wire capacitance. We assume that  $e\alpha V_g, g\mu_B B \ll E_{Fw}$ , with  $E_{Fw}$  the Fermi energy of the wire,  $g$  the Landé factor,  $\mu_B$  the Bohr magneton and  $e > 0$  the electronic charge. Then, the propagation of electrons with energy  $E \simeq E_{Fw}$  through the wire is described by a scattering matrix  $S_w = P(\theta_B) (\exp[i\delta\sigma^0 - i\sigma^z(\gamma_B/2)] \otimes \tau_1) P^{-1}(\theta_B)$  with  $\delta = L(k_{Fw} + [(E + e\alpha V_g - E_{Fw})/\hbar v_{Fw}])$ ,  $\gamma_B = g\mu_B BL/\hbar v_{Fw}$ ,  $P[\vartheta] = [\cos(\vartheta/2)\sigma_0 - i\sin(\vartheta/2)\sigma_y] \otimes \tau_0$ , and  $k_{Fw}$  ( $v_{Fw}$ ) the Fermi wave-vector (velocity) in the wire. Here,  $\sigma_i$ , with  $i \in \{0, x, y, z\}$ , are the Pauli matrices acting on spin space  $\{\uparrow_1, \downarrow_1\}$ . The matrices  $\tau_i$ , with  $i \in \{0, 1, 2, 3\}$ , are the Pauli matrices relating the space of incoming electrons  $\{(L, +), (R, -)\}$  to the space of outgoing electrons  $\{(L, -), (R, +)\}$ . We assume that there is no spin flip between the states  $\uparrow_n$  and  $\downarrow_n$  while electrons tunnel through interface  $n$ . Then, the scattering matrices describing the contacts 1 and 2 are respectively  $S_1 = \tilde{S}_1$  and  $S_2 = P[\theta]\tilde{S}_2P^{-1}[\theta]$  with  $2\tilde{S}_n = (\sigma_0 + \sigma_z) \otimes \tilde{s}_n^u + (\sigma_0 - \sigma_z) \otimes \tilde{s}_n^d$  and

$$\tilde{s}_n^s = \begin{bmatrix} r_{n,L}^s & t_{n,R}^s \\ t_{n,L}^s & r_{n,R}^s \end{bmatrix} \quad (1)$$

for  $s \in \{u, d\}$ . Here,  $t_{n,m}^s$  and  $r_{n,m}^s$  are complex amplitudes of transmission and reflection for electrons with spin  $s$ , incident from the side  $m \in \{L, R\}$  of contact  $n \in \{1, 2\}$ . We also define the transmission probability  $T_n^s = |t_{n,L(R)}^s|^2 = 1 - |r_{n,L(R)}^s|^2$ . We assume that electron-electron interactions can be neglected. Then, the linear conductance of the wire at temperature  $T$  is  $G = G_Q \sum_{s \in \{\uparrow_1, \downarrow_1\}, r \in \mathbb{B}} \int_0^{+\infty} \mathbb{T}_{sr}(E) (-\partial n_F(E)/\partial E)$ , with  $G_Q = e^2/h$ ,  $n_F(E) = 1/[1 + \exp(E/k_B T)]$  and  $\mathbb{T}_{sr}$  the probability that an electron with spin  $s$  from lead 1 is transmitted as an electron with spin  $r$  to lead 2 (we will use  $\mathbb{B} = \{\uparrow_1, \downarrow_1\}$  or  $\mathbb{B} = \{\uparrow_2, \downarrow_2\}$ , depending on convenience, for describing the spin state  $r$  in lead 2). The transmissions  $\mathbb{T}_{sr}$  can be found from the global scattering matrix  $S_{tot} = S_1 \circ S_w \circ S_2$  associated to the device

(see e.g. Ref. [17] for the definition of the composition rule  $\circ$ ). In the configurations studied in this Letter, the only interfacial scattering phases remaining in  $\mathbb{T}_{sr}$  are the reflection phases at the side of the wire, i.e.  $\varphi_1^{u[d]} = \arg(r_{1,R}^{u[d]})$  and  $\varphi_2^{u[d]} = \arg(r_{2,L}^{u[d]})$ . Importantly, these phases depend on spin because, due to the ferromagnetic exchange field, electrons are affected by a spin-dependent scattering potential when they encounter the interface between the wire and lead  $n$ . We will characterize this spin-dependence with the SDIPS parameters  $\Delta\varphi_n = \varphi_n^u - \varphi_n^d$ , with  $n \in \{1, 2\}$ . We also define the average tunneling rate  $T_n = (T_n^u + T_n^d)/2$  and the tunneling rate polarization  $P_n = (T_n^u - T_n^d)/(T_n^u + T_n^d)$ . In principle, the parameters  $T_n$ ,  $P_n$  and  $\Delta\varphi_n$  depend on the microscopic details of barrier  $n$ , but general trends can already be found from simple barrier models (see EPAPS Document for a detailed example). When there is a spin-independent barrier between the wire and lead  $n$ ,  $\Delta\varphi_n$  can be finite for  $T_n$  large only because too strong barriers prevent electrons from being affected by the lead properties. However, it is likely that the barrier between the wire and the lead is itself spin-polarized. This can be due to the magnetic properties of the contact material, but it can also be obtained artificially by using a magnetic insulator like e.g. EuS (see [18]) to form the barrier. In this case a large  $\Delta\varphi_n$  can be obtained for  $T_n$  small also. It is thus relevant to study the effects of the SDIPS (i.e. having  $\Delta\varphi_n \neq 0$ ) for a wide range of  $T_n$ .

We now present the results given by the scattering description of the circuit. We assume temperature  $T = 0$  and postpone a discussion on the effects of finite temperatures to the end of the paper. We first consider the case of parallel ( $\theta = 0$ , noted  $P$ ) or antiparallel ( $\theta = \pi$ , noted  $AP$ ) lead polarizations, with  $\theta_B = 0$ . We note  $\overline{\uparrow}_n = \downarrow_n$  and  $\overline{\downarrow}_n = \uparrow_n$ . In configuration  $c \in \{P, AP\}$ , one has  $\mathbb{T}_{ss}^c = 0$ , which means that spin is conserved when electrons cross the wire. The conductance of the device can be calculated from  $\mathbb{T}_{ss}^c = A_{ss}^c/|\beta_{ss}^c|^2$ , with  $s \in \{\uparrow_1, \downarrow_1\}$ ,  $A_{sr}^{P[AP]} = T_1^s T_2^{r[\overline{r}]}$  and  $B_{sr}^{P[AP]} = [(1 - T_1^s)(1 - T_2^{r[\overline{r}]})]^{1/2}$ . The term

$$\beta_{sr}^{P[AP]} = 1 - B_{sr}^{P[AP]} e^{i(\varphi_1^s + \varphi_2^{r[\overline{r}]}) + 2\delta + \kappa_s^1 \gamma_B} \quad (2)$$

with  $\kappa_{u(d)}^n = \pm 1$  for  $n \in \{1, 2\}$ , accounts for multiple reflections between the two contacts (we have used indices  $r \neq s$  and  $n \in \{1, 2\}$  in the above formulas for later use). The transmission probability  $\mathbb{T}_{ss}^{P[AP]}$  for spins  $s \in \{\uparrow_1, \downarrow_1\}$  is maximum at resonant energies  $E_s^{P[AP],j} = (2\pi j - \varphi_1^s - \varphi_2^{s[\overline{s}]} - \kappa_s^1 \gamma_B)(\hbar v_{Fw}/2L) - e\alpha V_g - E_{Fw}$ , with  $j \in \mathbb{Z}$ . Importantly, these resonant energies depend on spin  $s$  due to the SDIPS. This leads to the conclusion that the SDIPS can modify the conductance of a normal spin valve *even in a collinear configuration*. This feature is due to the ballistic nature of the system which offers the possibility of coherent multiple reflections between the contacts. Note that from Eq. (2), the SDIPS affects

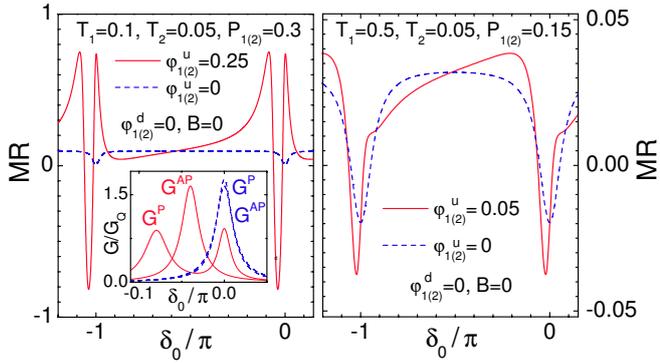


FIG. 2: Left panel: wire linear conductances  $G^P = G(\theta = 0)$  and  $G^{AP} = G(\theta = \pi)$  (inset) and magnetoresistance  $MR = (G^P - G^{AP})/(G^P + G^{AP})$  (main frame), as a function of the spin-averaged phase  $\delta_0$  acquired by electrons upon crossing the wire ( $\delta_0$  is linear with  $V_g$  in the limit studied in this paper). We used  $T_1 = 0.1$ ,  $T_2 = 0.05$ ,  $P_1 = P_2 = 0.3$ ,  $B = 0$ , and no SDIPS (blue dashed lines) or a SDIPS characterized by  $\varphi_1^u = \varphi_2^u = 0.25$  and  $\varphi_1^d = \varphi_2^d = 0$  (red full lines). The SDIPS produces a spin splitting of the conductance peaks because it shifts the phases accumulated by spins  $\uparrow$  and  $\downarrow$  upon reflection at the boundaries of the wire. This effect strongly increases the magnetoresistance of the device. Right panel: Magnetoresistance  $MR$  as a function of  $\delta_0$ , for a device with  $T_1 = 0.5$ ,  $T_2 = 0.05$ ,  $P_1 = P_2 = 0.15$ ,  $B = 0$  and no SDIPS (blue dashed lines) or a SDIPS such that  $\varphi_1^u = \varphi_2^u = 0.05$  and  $\varphi_1^d = \varphi_2^d = 0$  (red full lines). In this case, the spin-splitting of the conductance curves cannot be resolved but the symmetry of the MR oscillations is broken due to a conjugated effect of the SDIPS and the finite polarization of the transmission probabilities (see text).

electrons in the same way as a magnetic field collinear to the lead polarizations. However, the spin-splitting induced by the SDIPS can be different in the  $P$  and the  $AP$  configurations, contrarily to the splitting produced by a magnetic field collinear to  $\vec{p}_1$  and  $\vec{p}_2$ . Indeed, in the  $P[AP]$  configuration, there is a spin-splitting of the resonant energies if  $\Delta\varphi^{P[AP]} = \Delta\varphi_1 + [-]\Delta\varphi_2 \neq 0$ . In particular, the spin-splitting vanishes in the  $AP$  case when the contacts are perfectly symmetric. In the limit of low transmissions  $T_n^{u(d)} \ll 1$ , one can expand  $\mathbb{T}_{ss}^c(E)$  around  $E = E_s^{c,j}$  (see [16]) to obtain the Breit-Wigner formula  $\mathbb{T}_{ss}^c = A_{ss}^c [(2L[E - E_s^{c,j}]/\hbar v_{Fw})^2 + (T_1^s + T_2^{s[\bar{s}]})^2/4]^{-1}$ . This equation shows that the spin-splitting  $\Delta\varphi^c$  can be fully resolved in the conductance curve  $G^c(V_g)$  associated to configuration  $c$  if  $|\Delta\varphi^c| \gtrsim T_1^s + T_2^{s[\bar{s}]}$ , which we think possible in practice by using e.g. ferromagnetic insulators to make the contacts between the leads and the wire (see EPAPS Document). Figure 2, left panel, shows the conductances  $G^P$ ,  $G^{AP}$  and the magnetoresistance  $MR = (G^P - G^{AP})/(G^P + G^{AP})$  for a device with weakly transmitting barriers and  $B = 0$ , in the absence of SDIPS (blue dashed lines) or with a strong SDIPS such that  $\varphi_1^s = \varphi_2^s$  for  $s \in \{u, d\}$  (red full lines). For

convenience we plot the physical quantities as a function of  $\delta_0 = L(k_{Fw} + [e\alpha V_g/\hbar v_{Fw}])$  instead of the gate voltage  $V_g$ . The conductance presents resonances with a  $\pi$ -periodicity in  $\delta_0$ . As explained above, the SDIPS produces a spin-splitting of these resonances. Interestingly, this increases significantly the  $MR$  of the device by shifting the conductance peaks in the  $P$  and  $AP$  configurations. When  $\varphi_1^s \neq \varphi_2^s$ , both the  $G^P$  and  $G^{AP}$  curves can be spin-split, thus the  $MR$  curve can become more complicated (not shown), but this property persists as long as  $\Delta\varphi_1$  and  $\Delta\varphi_2$  remain larger than the transmission probabilities. When the transmissions become too large, it is not possible to resolve  $\Delta\varphi^c$  anymore because the dwell time of electrons on the wire decreases. Then, it is not possible to have a giant magnetoresistance. However, even in this situation, the SDIPS can modify qualitatively the  $MR$  of the device. Indeed, when there is no SDIPS, from the expression of  $\mathbb{T}_{ss}^c$ , the  $MR$  oscillations are always symmetric with  $V_g$ . Even a weak SDIPS can break this symmetry (see Fig. 2, right panel). The reason for this phenomenon is that although the transmission peaks associated to spins  $\uparrow$  and  $\downarrow$  are merged, these peaks have spin-dependent widths due to the polarizations  $P_{1(2)}$  of the transmissions. Thus, the position of the global maximum corresponding to  $E_{\uparrow}^{c,j}$  and  $E_{\downarrow}^{c,j}$  is different in the  $P$  and  $AP$  configurations. Interestingly, this allows to obtain  $MR(V_g)$  curves strikingly similar to the  $MR(V_g)$  measured by [19] in nanotubes connected to ferromagnetic leads [20].

We now study non-collinear configurations. When  $\theta \neq 0[\pi]$  and  $B = 0$ , one has, for  $s \in \{\uparrow_1, \downarrow_1\}$  and  $r \in \{\uparrow_2, \downarrow_2\}$

$$\mathbb{T}_{sr} = \frac{A_{sr}^P}{\left| \beta_{sr}^P \cos(\theta_{sr}/2) \left( 1 + \gamma_\phi^{\kappa_s^1 \kappa_r^2} \tan^2(\theta_{sr}/2) \right) \right|^2} \quad (3)$$

with  $\theta_{sr} = \widehat{(s, r)}$  and  $\gamma_\phi = \beta_{\uparrow_1 \downarrow_2}^P \beta_{\downarrow_1 \uparrow_2}^P / \beta_{\uparrow_1 \uparrow_2}^P \beta_{\downarrow_1 \downarrow_2}^P$ . The top left panel of Fig. 3 illustrates that the spin-splitting in  $G(\delta_0)$  goes continuously from  $|\Delta\varphi^P|/2$  to  $|\Delta\varphi^{AP}|/2$  when  $\theta$  goes from 0 to  $\pi$ . This can be used to tune the spin-splitting of the wire electronic spectrum. In the case  $c \in \{P, AP\}$  and  $\theta_B = \pi/2$ , one has, for  $(s, r) \in \{\uparrow_1, \downarrow_1\}^2$ , an expression analogue to Eq. (3), with  $\theta_{sr}$  replaced by  $\gamma_B$ ,  $\beta_{sr}^P$  replaced by  $\tilde{\beta}_{sr}^c$  with  $\tilde{\beta}_{sr}^{P[AP]} = 1 - \kappa_s^1 \kappa_r^1 e^{i(\varphi_1^s + \varphi_2^{r(\bar{r})} + 2\delta_0)}$ ,  $\gamma_\phi$  replaced by  $\tilde{\gamma}_\phi^c = \tilde{\beta}_{\uparrow_1 \downarrow_2}^c \tilde{\beta}_{\downarrow_1 \uparrow_2}^c / \tilde{\beta}_{\uparrow_1 \uparrow_2}^c \tilde{\beta}_{\downarrow_1 \downarrow_2}^c$  and  $\kappa_r^2$  replaced by  $\kappa_r^1$ . The signs  $\kappa_{s(r)}^1$  in  $\tilde{\beta}_{sr}^c$  account for the  $\pi$  phase shift acquired by a spin  $1/2$  when the magnetic field makes this spin precess by  $2\pi$ . Due to these signs, the dependence of  $G$  on  $\gamma_B$  can be very different from its dependence on  $\theta$ . For instance, in Fig. 3 (top panels),  $G(\delta_0, \gamma_B = 0)$  shows resonances close to  $\delta_0 = 0[\pi]$  for any value of  $\theta$  whereas the positions of the resonances in  $G^{P(AP)}(\delta_0)$  strongly vary with  $\gamma_B$ , with avoided crossings at  $\delta_0 \sim 0[\pi]$  in the  $P$  configuration and  $\delta_0 \sim \pi/2[\pi]$  in the  $AP$  configuration. We have already shown above that for  $\gamma_B = 0$ ,

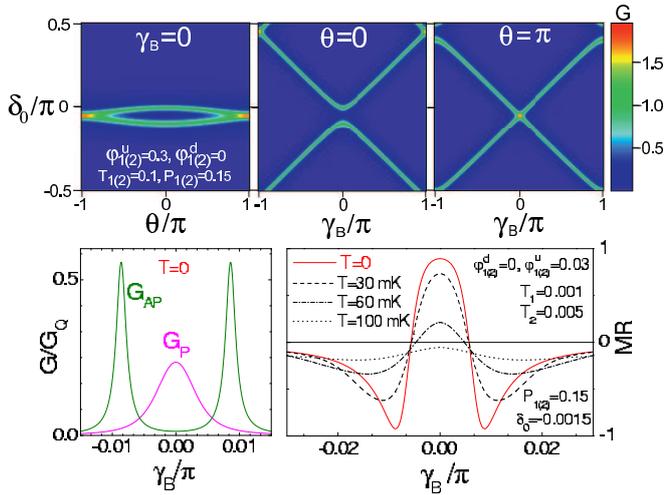


FIG. 3: Top: Color plots of the conductance  $G$  of the wire versus  $\delta_0$  and  $\theta$  for  $\gamma_B = 0$  (left panel) or versus  $\delta_0$  and  $\gamma_B$  for  $\theta = 0$  (middle panel) or  $\theta = \pi$  (right panel). We used  $T_1 = T_2 = 0.1$ ,  $P_1 = P_2 = 0.15$ ,  $\varphi_1^u = \varphi_2^u = 0.3$  and  $\varphi_1^d = \varphi_2^d = 0$ . Bottom: conductances  $G^P = G(\theta = 0)$  and  $G^{AP} = G(\theta = \pi)$  (bottom left panel) and magnetoresistance  $MR$  (bottom right panel) as a function of  $\gamma_B$ . We used  $T_1 = 0.001$ ,  $T_2 = 0.005$ ,  $P_1 = P_2 = 0.15$ ,  $\delta_0 = -0.0015$ ,  $\varphi_1^d = \varphi_2^d = 0$ , and  $\varphi_1^u = \varphi_2^u = 0.03$ . All the data are shown for temperature  $T = 0$  except the magnetoresistance which is shown also for finite temperatures by assuming that the wire has a length  $L = 500$  nm and a Fermi energy  $v_{Fw} \sim 8 \cdot 10^5$  m.s $^{-1}$ . The  $MR$  changes sign abruptly at a low value of  $\gamma_B$  because the resonances in  $G^{AP}(\gamma_B)$  and  $G^P(\gamma_B)$  are shifted for the value of  $\delta_0$  considered (this configuration is possible thanks to the spin-dependent resonance pattern, as can be understood from the color plots). The effect persists up to 90 mK for the parameters chosen here. In this figure, we used  $\varphi_1^s = \varphi_2^s$ , thus there is no spin-splitting of the resonant energy at ( $\delta_0 = 0, \theta = \pi$ ). However, we have checked that the field effect shown in the bottom panel can persist at  $\varphi_1^s \neq \varphi_2^s$ .

electrons are affected by an SDIPS-induced effective-field which depends on the electrode configuration. As a consequence, for low transmissions, it is possible to find values of  $\delta_0$  such that the conductances  $G^P$  and  $G^{AP}$  versus  $\gamma_B$  display distinct resonances close to  $\gamma_B = 0$  (see bottom left panel of Fig. 3). This allows to obtain a giant magnetoresistance whose sign can be switched by applying a magnetic field with  $\gamma_B \ll 1$  (bottom right panel of Fig. 3). For instance, for  $T_1 = 0.001$ ,  $T_2 = 0.005$ ,  $P_n = 0.15$ ,  $\Delta\varphi_n = 0.03$ ,  $\delta_0 = -0.0015$ ,  $L = 500$  nm and  $v_{Fw} \sim 8 \cdot 10^5$  m.s $^{-1}$  [14], one has  $MR \sim +89\%$  at  $B = 0$  and  $MR \sim -92\%$  at  $B = 250$  mT (Fig. 3). This small value of magnetic field is particularly interesting since in practice, it is difficult to keep  $\vec{p}_1$  and  $\vec{p}_2$  perpendicular to  $\vec{B}$  when  $B$  becomes larger than typically 1 T.

We now briefly address the effect of finite temperature  $T$ , which starts to modify the behavior of the circuit when it becomes comparable to the wire energy-level spacing  $\hbar v_{Fw}/2L$  times the transmission probabilities (see the

above Breit-Wigner formula). The switching of the  $MR$  sign with  $V_g$  described in Fig. 2, left, is relatively robust to temperature since it is still obtainable at 1 K for the wire parameters considered in the previous paragraph (not shown). For Fig. 3, having a switching of the  $MR$  sign with a low magnetic field requires to have lower temperatures due to the low values of transmission probabilities necessary. However, this effect should be obtainable in practice since it persists up to 90 mK for the wire considered here (see Fig. 3).

So far, we have disregarded the gate-dependence of the scattering matrices  $S_{1(2)}$ . This is correct if the variations of  $e\alpha V_g$  are negligible compared to the characteristic energy scales defining the interface scattering potentials. However, the opposite situation can occur (see EPAPS Document for an example). In this case, the effective field produced by the SDIPS is gate-dependent. The most simple consequence of this feature is that the SDIPS-induced spin-splitting of the resonant energies depends on the resonance index  $j$ . The gate-dependence of the SDIPS effective field could be used for controlling the spin dynamics.

In summary, we have studied the effects of the spin-dependence of interfacial phase shifts (SDIPS) on the linear conductance of a ballistic one-dimensional quantum wire connected to two ferromagnetic leads. The SDIPS generates a spin-splitting of the wire energy spectrum which is tunable with the gate voltage and the angle between the ferromagnetic polarizations. This leads in particular to a giant magnetoresistance effect with a sign tunable with the gate voltage and the magnetic field. These properties could be exploited for manipulating spins in the context of spin electronics or quantum computing.

This work was supported by the RTN Spintronics, the Swiss NSF, the RTN DIENOW, and the NCCR Nanoscience.

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- [20] Ref. [19] shows data in terms of  $MR' = (G^P - G^{AP})/G^{AP}$ , but for low polarizations, one has  $MR' \sim 2MR$ . A more quantitative interpretation of these data would require to take into account Coulomb blockade effects observable in this experiment, but we expect that the SDIPS-induced spin-splitting persist in this situation.

# Auxiliary material for the Letter “Controlling spin in an electronic interferometer with spin-active interfaces”

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(Dated: November 29, 2005)

We evaluate the parameters  $T_n$ ,  $P_n$ , and  $\Delta\varphi_n$  for the case of a Dirac potential barrier between the wire and lead  $n$ . This study indicates that it is relevant to study the effects of the SDIPS (i.e., having  $\Delta\varphi_n \neq 0$ ) for a wide range of  $T_n$ . We also find that  $\Delta\varphi_n$  can depend on the gate voltage  $V_g$ , which would be useful for tuning the SDIPS-induced spin-splitting.

In principle, the parameters  $T_n$ ,  $P_n$ , and  $\Delta\varphi_n$  describing the scattering of electrons at barrier  $n$  depend on the microscopic details of this barrier. However, it is interesting to study a Dirac barrier model which already gives important general trends [1]. In this model, one has, for  $s \in \{u, d\}$ ,  $r_{1,R(2,L)}^s = (k_{Fw} - k_{1(2)}^s - i[2m_e U_{1(2)}^s/\hbar^2]) / (k_{Fw} + k_{1(2)}^s + i[2m_e U_{1(2)}^s/\hbar^2])$ . Here,  $k_{Fw}$  and  $k_n^s$  are the Fermi electronic wave-vectors for spin  $s$ , in the wire and in lead  $n$  respectively and  $U_n^s$  is the strength of the Dirac barrier for spin  $s$ . The modulus of the polarization  $\vec{p}_n$  of lead  $n$  is given by  $p_n = (N_{F,n}^u - N_{F,n}^d) / (N_{F,n}^u + N_{F,n}^d)$ , with  $N_{F,n}^s$  the density of states at the Fermi level for spin  $s$  in lead  $n$  ( $N_{F,n}^u > N_{F,n}^d$ ). Then, in an s-band Stoner model, one has  $(\hbar k_n^{u[d]})^2 / 2m_e E_F = 1 + [-]2p_n / (1 + p_n^2)$ , with  $m_e$  the free electron mass. In Fig. 1, we show the average tunneling rate  $T_n = (T_n^u + T_n^d)/2$ , the tunneling rate polarization  $P_n = (T_n^u - T_n^d) / (T_n^u + T_n^d)$  and the SDIPS parameter  $\Delta\varphi_n$  predicted with this approach. We have used the value  $k_{Fw} = 8.5 \cdot 10^9 \text{m}^{-1}$  typical of single-wall nanotubes [2], and the Fermi energy  $E_F = 10 \text{eV}$  for lead  $n$ . We have also assumed  $V_g = 0$  and  $B = 0$ . When the barrier is considered to be spin-independent, i.e.  $\alpha_n = (U_n^\downarrow - U_n^\uparrow) / (U_n^\uparrow + U_n^\downarrow) = 0$ ,  $\Delta\varphi_n$  can be finite for  $T_n$  large only. Figure 1 also shows the case  $\alpha_n > 0$  i.e., the barrier is also magnetically polarized, with the same polarization direction as the ferromagnet. This can be caused by the magnetic properties of the contact material, but it can also be obtained artificially by using a magnetic insulator to form the barrier. In this case a large  $\Delta\varphi_n$  can be obtained for  $T_n$  small also. It is thus relevant to study the effect of the SDIPS (i.e. having  $\Delta\varphi_n \neq 0$ ) for a wide range of  $T_n$ .

In the main body of the Letter, we disregarded the gate-dependence of the scattering matrices  $S_{1(2)}$ . However, the validity of this approximation depends on the characteristics of the circuit. As an example, we consider a wire with  $v_{Fw} \sim 8 \cdot 10^5 \text{m.s}^{-1}$  and  $L = 500 \text{nm}$ , connected to two barriers like that described by the full lines in Fig. 1 of this EPAPS document, with  $T_n \sim 0.1$ . The oscillation period in  $G(V_g)$  is  $T_g = \hbar v_{Fw} \pi / e\alpha L$ . Starting

from  $V_g = 0$  for which  $\Delta\varphi^P \sim 0.4$ ,  $\Delta\varphi^P$  varies by only 0.15% when  $V_g$  changes by  $2T_g$ . Thus, on this scale, the SDIPS can be considered as constant with  $V_g$  and the previous approach is correct. On larger scales, the periodicity of the  $G(V_g)$  curves is broken and the SDIPS-induced spin-splitting of the resonant energies can be tuned with  $V_g$ . For instance, a shift of  $\delta_0$  by  $70T_g$  makes  $\Delta\varphi^P$  vary by  $\sim 5\%$ , i.e.,  $\sim 100 \text{mT}$  in terms of effective magnetic field [3]. This could be used for controlling the spin dynamics.

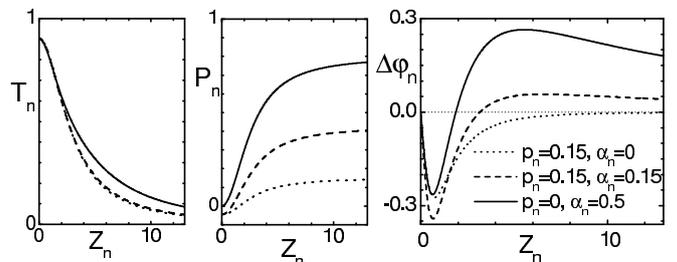


FIG. 1: Spin-averaged tunneling rate  $T_n$  (left panel), tunneling rate polarization  $P_n$  (middle panel) and SDIPS parameter  $\Delta\varphi_n$  (right panel) estimated by modeling the contact  $n$  with a Dirac barrier with coefficient  $U_n^s$ , placed between a ferromagnetic metal with Fermi energy  $E_F = 10 \text{eV}$  and a wire with  $k_{Fw} = 8.5 \cdot 10^9 \text{m}^{-1}$  [2]. We show the results as a function of the average barrier impedance  $Z_n = m_e(U_n^u + U_n^d) / (\hbar^2 k_{Fw})$ , for different values of the polarization  $p_n$  of lead  $n$  and of the spin asymmetry  $\alpha_n = (U_n^\downarrow - U_n^\uparrow) / (U_n^\uparrow + U_n^\downarrow)$  of the barrier. In the case  $\alpha_n > 0$  and  $Z_n$  large, one has  $\Delta\varphi_n > 0$  due to a weaker penetration of minority electrons in the barrier. When  $\alpha_n = 0$  or  $Z_n$  small, the wave-vector mismatch between the lead and the wire leads to  $\Delta\varphi_n < 0$  (this mismatch can also lead to  $P_n < 0$  for  $Z_n$  small).

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[3] This can be done while fulfilling the assumption  $e\alpha V_g \ll E_{Fw}$  made in the article.