Tunable Conductance of Magnetic Nanowires with Structured Domain Walls

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We show that in a magnetic nanowire with double magnetic domain walls, quantum interference results in spin-split quasistationary states localized mainly between the domain walls. Spin-flip-assisted transmission through the domain structure increases strongly when these size-quantized states are tuned on resonance with the Fermi energy, e.g., upon varying the distance between the domain walls which results in resonance-type peaks of the wire conductance. This novel phenomenon is shown to be utilizable to manipulate the spin density in the domain vicinity. The domain wall parameters are readily controllable, and the predicted effect is hence exploitable in spintronic devices.

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The discovery of the giant magnetoresistance [1] and its rapid and diverse industrial utilizations sparked major efforts in understanding and exploiting spin-dependent transport phenomena. In particular, new perspectives of even broader importance are anticipated from combining semiconductor technology with nanoscale fabrication techniques to produce magnetic semiconducting materials and control precisely the spins of the carriers. At the heart of this new field that is now termed ''semiconductor spintronics'' [2] is the understanding of the transport properties of a magnetic domain wall (DW), which is a region of inhomogeneous magnetization between two domains of homogeneous (different) magnetizations. Thick (or adiabatic) DWs occurring in bulk metallic ferromagnets have an extension much larger than the carriers' Fermi wavelength and are largely irrelevant for the resistance [3]. In contrast, a series of recent experiments on magnetic nanostructures, and particularly nanowires, revealed that the magnetoresistance in the presence of DWs can be as large as several hundreds [4,5] or even thousands [6,7] of percents. These observations have decisive consequences in so far as DWs are controllable efficiently by applying a magnetic field [8] and can also be steered by a spin-polarized electric current [9], meaning that the magnetoresistance of the structure containing DW is controllable via an electric field.

The interpretation of the huge magnetoresistance of DWs observed in the magnetic semiconductor (in the ballistic quantum regime) relies on the relative sharpness of DWs on the scale set by the wavelength of carriers (electrons or holes) [10–14]. In such a situation spin-dependent scattering of carriers from DWs is greatly enhanced. In this Letter we predict the formation of spin quantum wells and the occurrence of a novel effect in magnetic semiconductor nanowires with double DWs: By pinning one of the DWs at a constriction, one can control the location of the second DW. The spin-dependent transmission and reflection of carriers waves from the first and the second DWs and the quantum interference between these waves lead to the formation of spin-split quasistationary states, and hence the double DWs act in effect as a penetrable ''spin quantum well'' (located between the two DWs). Lifetimes and energetic positions of these states depend on experimentally controllable parameters of DWs such as the extent of the well, i.e., on the distance between DWs. As a result, the conductance of a structure with DWs possesses sharp resonances when the Fermi energy matches the spin quantum-well states. Hence, the resistance varies by several orders in magnitude in response to minor changes in DW positions or in carrier density, e.g., achieved upon gating the structure. Strong backscattering from DWs and interferences between carrier waves lead to spin-density accumulation that can be tuned externally by modifying the DWs structure.

We consider a magnetic wire with a magnetization profile exhibiting two DWs separated by the distance $2L$. The magnetization vector field $\mathbf{M}(z)$ in both DWs varies within the x-z plane. The z axis is along the wire. Thus, z is the easy axis, and the x-z plane is the easy plane. We study the case where the thickness and the width of the wire are smaller than the carrier Fermi wavelength so that only one size-quantized level (a single one-dimensional subband) is populated. This situation is realizable in magnetic semiconductor-based structures. The Hamiltonian describing independent carriers along the wire in the presence of the magnetization field $\mathbf{M}(z)$ is

$$H = -\frac{\hbar^2}{2m}\frac{d^2}{dz^2} - JM_z(z)\sigma_z - JM_x(z)\sigma_x, \qquad (1)$$

where J is the exchange coupling constant, $M_{x(z)}$ is the $x(z)$ component of the inhomogeneous magnetization field, and m is the carrier's effective mass. Figure 1 shows a schematic drawing of the spin-up and spin-down band-edge profiles that are utilized for a full quantal treatment of the spin-dependent scattering of charge carriers.

We are interested in the case where the width 2δ of each DW is smaller than the carriers Fermi wavelength, $k_F\delta \ll 1$ and particularly when $k_F L \geq 1$ (otherwise the carriers are not influenced by the detailed topology of the DWs). For moderate carrier density the chemical potential μ is in one of the magnetically split subbands (cf. Fig. 1). This case corresponds to a full spin polarization of the electron gas, a situation realizable in magnetic semiconductor nanowires. Carrier wave functions are expressible as

$$\psi_k(z) = (e^{ikz} + re^{-ikz})|\uparrow\rangle + r_f e^{\kappa z}|\downarrow\rangle, \qquad z < -L, \quad (2)$$

$$\psi_k(z) = (Ae^{\kappa z} + Be^{-\kappa z})|\uparrow\rangle + (Ce^{ikz} + De^{-ikz})|\downarrow\rangle, \qquad |z| < L, \quad (3)$$

$$\psi_k(z) = te^{ikz}|\uparrow\rangle + t_f e^{-\kappa z}|\downarrow\rangle, \qquad z > L, \quad (4)$$

where $|\uparrow\rangle$ ($|\downarrow\rangle$) is the spin-up (spin-down) component of the carrier states, $k = [2m(\varepsilon + JM)]^{1/2}/\hbar$, and $\kappa = [2m(JM - \varepsilon)]^{1/2}/\hbar$. The electron energy ε is measured from the midpoint between spin-up and spin-down band edges. The non-spin-flip (spin-flip) transmission and reflection coefficients t and r (t_f and r_f) as well as the constants A, B, C, and D have to be deduced from the solutions of Eq. (1) and from wave function continuity requirements. Physically, Eqs. (2)–(4) describe spin-up carriers incoming from the left, being transmitted and reflected from the double DW structure into waves with the same or opposite spin polarizations with a subsequent decay of the spin-down part of the wave function.

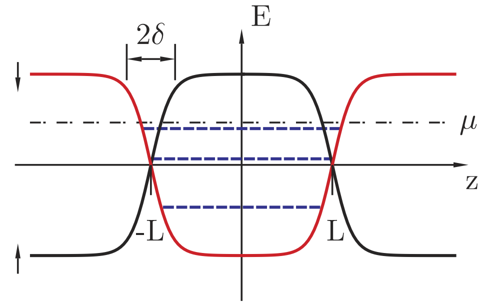

FIG. 1 (color online). Schematics of the double domain wall structure: indicated are the potential energy profiles for spin-up and spin-down electrons. Dashed lines between domain walls stand for quasilocalized energy levels in the spin-down quantum well.

To determine the unknown coefficients in Eqs. (2)–(4), we utilize the wave function continuity at $z = \pm L$, i.e.,

$$\frac{\hbar}{2m}\left(\frac{d\psi_k}{dz}\bigg|_{z=-L+\delta} - \frac{d\psi_k}{dz}\bigg|_{z=-L-\delta}\right) + \lambda\sigma_x\psi_k(-L) = 0, \quad (5)$$

where

$$\lambda \simeq \frac{J}{\hbar}\int_{-L-\delta}^{-L+\delta} dz M_x(z) \simeq \frac{2JM\delta}{\hbar}. \quad (6)$$

Similar equation holds for $z = L$. The boundary conditions at $z = \pm L$ (eight equations for the spinor components) define all the coefficients in Eqs. (2)–(4), e.g.,

$$[\kappa - ik\tanh(2\kappa L)]\tilde{r} - \Delta\tanh(2\kappa L)\tilde{r}_f - \frac{\kappa\tilde{t}}{\cosh(2\kappa L)} = -[\kappa + ik\tanh(2\kappa L)]e^{-ikL}, \quad (7)$$

$$-\Delta\sin(2kL)\tilde{r} + [\kappa\sin(2kL) + k\cos(2kL)]\tilde{r}_f - k\tilde{t}_f = \Delta\sin(2kL)e^{-ikL}, \quad (8)$$

$$-[\kappa\tanh(2\kappa L) - ik]\tilde{r} + \Delta\tilde{r}_f + \frac{ik\tilde{t}}{\cosh(2\kappa L)} + \frac{\Delta\tilde{t}_f}{\cosh(2\kappa L)} = [\kappa\tanh(2\kappa L) + ik]e^{-ikL}, \quad (9)$$

$$\Delta\cos(2kL)\tilde{r} - [\kappa\cos(2kL) - k\sin(2kL)]\tilde{r}_f + \Delta\tilde{t} - \kappa\tilde{t}_f = -\Delta\cos(2kL)e^{-ikL}. \quad (10)$$

Here we introduced the following notation: $\Delta = 2m\lambda/\hbar$, $\tilde{r} = re^{ikL}$, $\tilde{r}_f = r_f e^{-\kappa L}$, $\tilde{t} = te^{ikL}$, and $\tilde{t}_f = t_f e^{-\kappa L}$. In this way we derived (rather cumbersome) analytical expressions for the reflection and transmission coefficients.

For physical insight into the results we inspect at first the limiting situation $\Delta = 4mJM\delta/\hbar^2 = 0$ in which case no spin-flip transitions occur at the DW. Correspondingly, only spin-up electrons tunnel through the barrier. The standard formula for barrier tunneling $t = 2ik\kappa e^{-2ikL}[2ik\kappa\cosh(2\kappa L) + (k^2 - \kappa^2)\sinh(2\kappa L)]^{-1}$ is retrieved using Eqs. (7) and (8). The spin-down electrons are localized within the spin quantum well. From Eqs. (9) and (10) we find the symmetric solution (with $r_f = t_f$) that corresponds to localized states with the wave vector k and obeys the relation $\tan(kL) = \kappa/k$. The antisymmetric solution for such k (with $r_f = -t_f$) satisfies the equation $\tan(kL) = -\kappa/k$. When the distance L between the DWs is varied, the energetic positions of the size-quantized levels within the well are shifted. For certain values of L the energy of the localized states within the well (dashed lines in Fig. 1) coincide with the Fermi level. Thus, if the spin-mixing amplitude is finite (i.e., $\Delta \neq 0$), we expect spin-up carriers to transverse resonantly the DWs. It is important to note here that when $\Delta \neq 0$, the aforementioned localized spin-down states (dashed lines in Fig. 1) turn quasistationary with a finite decay width $\Gamma = \hbar/\tau \sim |t_f|^2$ where τ is the lifetime of these quasilocalized states. $|t_f|$ is controllable, e.g., by changing the parameters of the DWs.

In the regime of a linear response to an applied bias voltage the conductance of the wire is determined by the transmission coefficient t as

$$G = \frac{e^2}{2\pi\hbar}|t(\varepsilon = \mu)|^2. \quad (11)$$

Using this relation and the solutions derived from Eqs. (7)–(10), we calculated the variation of the conductance G with the DWs distance ($2L$) for several values of the magnetization M. The conductance shown in Fig. 2 exhibits narrow resonance peaks corresponding to those values of L at which a quasidiscrete, size-quantized level coincides with the Fermi energy. At the conductance peak position the effective barrier created by the DWs is basically transparent. The width of the resonance peaks is related to τ, the lifetime of the quasistationary, spin-well states and is determined by the spin-mixing mechanism mentioned above (and specifically by $|t_f|$). As demonstrated by Fig. 3, the strength of spin mixing and hence the width of the resonance peaks can be controlled, e.g., by varying the width δ of the DWs. Decreasing the spin-mixing parameter $\Delta = 4mJM\delta/\hbar^2$, the lifetime of the localized spin quantum-well states increases and the conductance resonance peaks become correspondingly narrower. The energetic positions of the quasidiscrete levels depend also on Δ (and hence on δ). This results in a slight shift of the resonance positions when changing δ, as shown in Fig. 3. Experimentally the Fermi level position can be shifted by electrically gating the whole structure. In this case, the resonance conductance peaks occur as a function of the gate voltage for a fixed distance between the DWs.

It is important to note that within our model the resonant transmission does not vanish in the limit of large L, for we do not incorporate decoherence effects that destroy the interference of transmitted and reflected waves in the region between the magnetic DWs. Thus, our present considerations are valid for $L \ll L_e$, where L_e is the decoherence length. On the other hand, we expect the double DWs resistance in this limit to be the sum of the resistances of the individual DWs that we calculated previously [14]. Furthermore, we note that, while our calculations are done for $T = 0$, the effect of the temperature T is negligibly small as long as $T \ll E_0$, where E_0 is the effective barrier for spin-up electrons, $E_0 = JM - \mu$. For higher temperatures, the activated spin-up electrons contribute to the conductance, and the resonant character of the conductance is smeared out.

A further notable feature of the scattering of spin-split carriers from DWs is the interference-induced buildup of spin density in the DWs vicinity, as demonstrated in Fig. 4. The period of the spin-density oscillation is determined by the Fermi wave vector [close to the DWs there is a fast-decaying contribution (proportional to $e^{\kappa z}$) that is superimposed on the simple oscillations shown in Fig. 4]. The degree of the spin-density accumulation can be controlled, to a certain extent, by changing experimentally the parameters of the DWs. For instance, for the well size $L \simeq 14.66$ nm quasistationary well states are formed and are energetically close to the Fermi energy. As a result, DWs are almost nonreflecting and the spin-density buildup diminishes (cf. Fig. 4). Changing L (i.e., off resonance) DWs backscatter strongly and large spin density is accumulated. By gating the structure, we can tune the Fermi energy and manipulate the spin density in a manner similar to Fig. 4. For an experimental verification we note that spin-density modulations can be imaged with a subnanometer resolution using spin-polarized scanning tunneling microscopy [15]. Hence, the predictions of Fig. 4 are accessible experimentally.

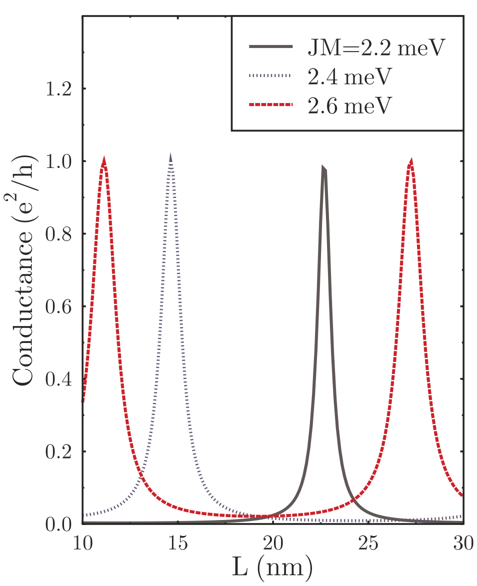

FIG. 2 (color online). Conductance vs distance between two domain walls for different values of the magnetization M. For the numerical calculations we assumed $m = 0.6m_0$ (m_0 is the free electron mass), $\delta = 2$ nm, and $\mu = -2$ meV.

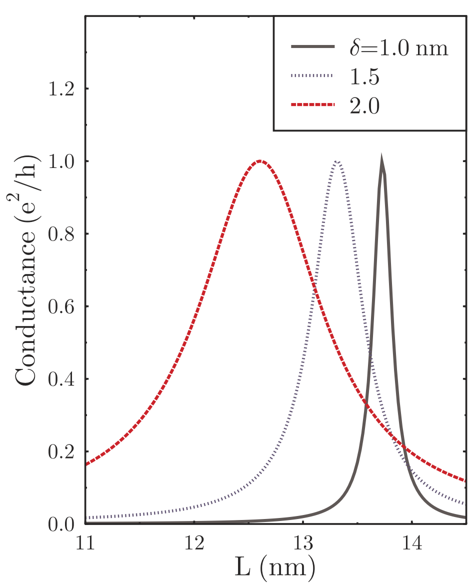

FIG. 3 (color online). Conductance vs distance between two domain walls for different values of the domain wall width δ. Other parameters are $JM = 2.5$ meV, $m = 0.6m_0$, and $\mu = -2$ meV.

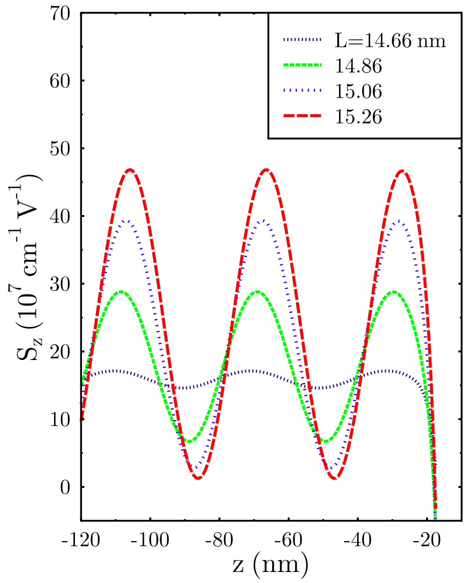

FIG. 4 (color online). Spin-density profile for $z < -L$ at different values of spin quantum well L. Here S_z is the density of electron spins per unit wire length and unit voltage bias. The parameters are $JM = 2.4$ meV, $\mu = -2$ meV, and $\delta = 2$ nm.

It is useful to compare the present findings with resonant tunneling for manipulating the spin orientation in magnetic layered structures [16] in which case the discrete levels are controlled via an external magnetic field that modifies the spin splitting of quantum-well energy levels. In contrast, our proposition relies on an interference-induced creation of a *spin quantum well* by the presence of double DW structure and as demonstrated above offers a range of external parameters with which the conductance can be tuned. For an experimental realization nanowires of low-density magnetic semiconductors [17] (5×10^{18} cm^{-3}) are favorable. Such structures are reported in [7]; however, the DWs separation was too large (500 nm) for a noticeable effect. In principle, however, the experiment should be feasible with the parameters employed above: e.g., the wire Fermi momentum is $k_F = \pi\rho_{1D}$, where ρ_{1D} is the linear hole density related to the bulk density by $\rho_{1D} = \rho_{3D}S$ and S is the wire cross section. For $\rho_{3D} = 10^{20}$ cm^{-3} and $S = 1$ nm^2, we obtain $k_F \simeq 3 \times 10^6$ cm^{-1} corresponding to a wavelength $\lambda_F \simeq 20$ nm and an energy $E_F \simeq 5$ meV. These numbers are in the range of those used in our calculations.

Our considerations assumed electrons as carriers. In III-V magnetic semiconductors the carriers are holes with an energy spectrum more complicated than that described by Eq. (1). It is clear, however, that in this case carrier scattering and inferences lead to the formation of a spin quantum well with localized levels, and hence the physical phenomena discussed above are expected to emerge as well for the case of hole carriers. For a strong magnetization and hence large splitting of the valence subbands, $JM \gg |\mu|$ (μ is the chemical potential measured from the valence band edge) one can employ a model with parabolic bands as for electrons.

To summarize, the conductance of a magnetic nanowire with double DWs is shown to possess a strong dependence on the DWs separation. The extreme sensitivity of the conductance on the interwalls' distance can be used to identify the relative position of the DWs. It can also be utilized to transform a magnetic field effect on the DWs into a change of the current flowing through the nanowire. A wall displacement of the order of 10% induces a resistance change of hundreds of percents.

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