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Materials Science and Engineering B 126 (2006) 254–257

MATERIALS SCIENCE & ENGINEERING B

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Resonant transmission through a double domain wall in magnetic nanowires

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Abstract

As experimentally verified, a large magnetoresistance arises due to domain walls creation (or destruction) in Ni nanowires and in some nanostructures based on GaMnAs magnetic semiconductors. Hence the presence and structuring of magnetic domain walls have important potential applications in magnetoelectronics devices. Here, we uncover a way of controlling the conductance via resonant transmission through a double domain wall structure. This phenomenon is due to quantum interference of charge carrier wave functions in spin quantum wells, which leads to the formation of quantized energy states in the potential well created by a double domain wall. When the energy of a state in the spin quantum well is resonant with the Fermi energy in the wire, the spin–flip transmission through the domain walls becomes most effective. This gives rise to a resonance-type dependence of the conductance on the distance between the domain walls or on the Fermi energy.

Keywords: Domain walls; Magnetoresistance; Magnetic semiconductors

1. Introduction

Recent experiments on the magnetoresistance (MR) of magnetic nanowires with domain walls (DW) have indicated that MR can be as large as several hundreds [1,2] to thousands [3,4] of percents. Theoretically, the mechanism behind this remarkable phenomenon is still not yet established. On the other hand, the theory of MR in the case of a smooth domain wall in a bulk ferromagnet (domain wall width much larger than the Fermi wavelength) has long been put forward and it predicts no appreciable influence of the domain wall on MR. This behaviour obviously changes drastically when the size of the domain shrinks, as demonstrated recently for Ni [5].

The situation changes in magnetic semiconductor-based structures. The Fermi wavelength of charge carriers may well exceed the size of the domain wall, and the conventional semiclassical theory of MR in the presence of domain walls is not applicable. Indeed, as observed recently these structures show a huge MR effect [4]. On the theoretical side several (analytical, and full-numerical) methods have been developed (see Refs. [6–11]) to account for the huge MR due to sharp domain walls in the regime of ballistic transport. These studies are fuelled not only by the conceptual and fundamental interest of DW MR, but also by possible applications in spintronics [12]. For example, properties of DWs are controllable by applying a weak magnetic field [13] or driving an electric current [14] through the device. Hence the DWs and the associated MR can be manipulated in a precise way by changing the external parameters.

This current study is focused on a novel effect, namely of the spin-dependent resonance-type phenomena in a magnetic nanowire with sharp double DWs. One of these two walls can be pinned by a constriction while the spatial position of the second DW can be controllably shifted. As demonstrated below, energy states of charge carriers in the spin quantum well created between the DWs are quantized. This quantization leads to a large (resonant) change in the MR of the structure upon a small change of the domain wall position (the latter change implies a shift of the energetic position of the resonant states in the spin quantum well).

2. Model of the double domain wall

We consider a quasi-one-dimensional magnetic wire with a magnetization profile consisting of two domain walls separated

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doi:10.1016/j.mseb.2005.09.024

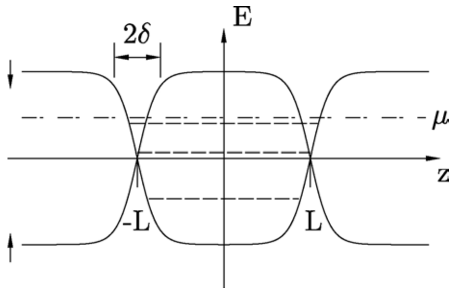

Fig. 1. A schematic picture of a double domain wall structure in a magnetic wire. Solid lines stand for the potential energy profiles for spin-up and -down carriers. Quasi-localized energy levels are shown by dashed lines in the spin-down quantum well between the domain walls.

by distance $2L$. For both DWs the magnetization vector $\mathbf{M}(z)$ is in the x–z plane. The z axis is chosen to be along the wire. So, the x–z plane is the easy plane of the ferromagnet. We assume that only one transverse channel is relevant, i.e. the transverse dimensions of the wire are small compared to the Fermi wavelength of charge carriers; a situation typically achieved in magnetic semiconductors.

The Hamiltonian H describing charge carriers interacting with the magnetization $\mathbf{M}(z)$ takes the form:

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

where J is the exchange coupling constant. Fig. 1 shows a schematic drawing of the spin-up and -down band-edge profiles. It is important to note that our treatment below is full quantum mechanical, even though the presentation of the energy band-edge is purely quasi-classical.

As stated above the situation of interest here is when the width of each of the DWs 2δ is much smaller than the wavelength of charge carriers at the Fermi level, $k_F\delta \ll 1$. In addition, we are particularly interested in effects occurring when $L \geq k_F^{-1}$. It is then justified to model DWs as infinitely thin ones (on the length scale of k_F^{-1}), with some spin–flip scattering at the walls [10].

We assume that the density of carriers is sufficiently low, so the chemical potential μ lies in one of the spin–split subbands, as depicted in Fig. 1. This situation corresponds to a full spin polarization of the electron gas, and it can be readily realized in magnetic semiconductor wires.

The scattering states of the charge carriers are described by the wave functions:

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

$$= (A\,e^{\kappa z} + B\,e^{-\kappa z})\,|\uparrow\rangle + (C\,e^{ikz} + D\,e^{-ikz})\,|\downarrow\rangle, \quad -L < z < L, \tag{3}$$

$$= t\,e^{ikz}\,|\uparrow\rangle + t_f\,e^{-\kappa z}\,|\downarrow\rangle, \qquad z > L, \tag{4}$$

where $|\uparrow\rangle$ and $|\uparrow\rangle$ are the spin-up and spin-down states, $k = [2m(\varepsilon + JM)]^{1/2}/\hbar$, $\kappa = [2m(JM - \varepsilon)]^{1/2}/\hbar$, t, t_f, r, and r_f are the transmission and reflection coefficients, respectively, A, B, C, and D are constants to be determined from appropriate boundary conditions, and M is the absolute value of the magnetization.

At $z = \pm L$ the wave function is continuous. However, its derivative suffers a jump, which at $z = -L$ is

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

where

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

Similar expressions also hold for $z = L$.

The conditions at the points $z = \pm L$ lead to eight equations for the spinor components and therefore define all the coefficients in Eqs. (2)–(4). The number of equations can be reduced by excluding the coefficients A, B, C, D, which correspond to the wave function in the regionbetween the domain walls. This leads to the following set of equations:

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

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

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

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

Here we use the shorthand notation: $\Delta = 2m\lambda/\hbar$, $\tilde{r} = r\,e^{ikL}$, $\tilde{r}_f = r_f\,e^{-\kappa L}$, $\tilde{t} = t\,e^{ikL}$, and $\tilde{t}_f = t_f\,e^{-\kappa L}$. The above equations can be solved analytically, but the corresponding formulas are rather cumbersome and will not be presented here.

In the case of $\Delta = 0$, no spin–flip transitions occur at the DW and hence only spin-up electrons can traverse the wall. In this situation we make use of Eqs. (7) and (8) and retrieve the standard formula for the the tunnelling amplitude:

$$t = \frac{2ik\kappa\,e^{-2ikL}}{2ik\kappa\,\cosh(2\kappa L) + (k^2 - \kappa^2)\,\sinh(2\kappa L)}. \tag{11}$$

The spin-down electrons are located in the spin quantum well. Using Eqs. (9) and (10) we derive the symmetric solution ($r_f = t_f$) which corresponds to the localized states with k obeying the equation $\tan(kL) = \kappa/k$. The antisymmetric solution ($r_f = -t_f$) occurs for k obeying $\tan(kL) = -\kappa/k$. When varying distance L between the DWs, the size-quantized energy levels in the well are shifted. At some values of L one of these levels coincides with the Fermi level. Thus, for a finite spin-mixing amplitude, $\Delta \neq 0$, one can expect a resonant transmission through the double DW.

3. Resonant transmission

In the linear response regime to an external voltage the conductance G of the structure is determined by the transmission

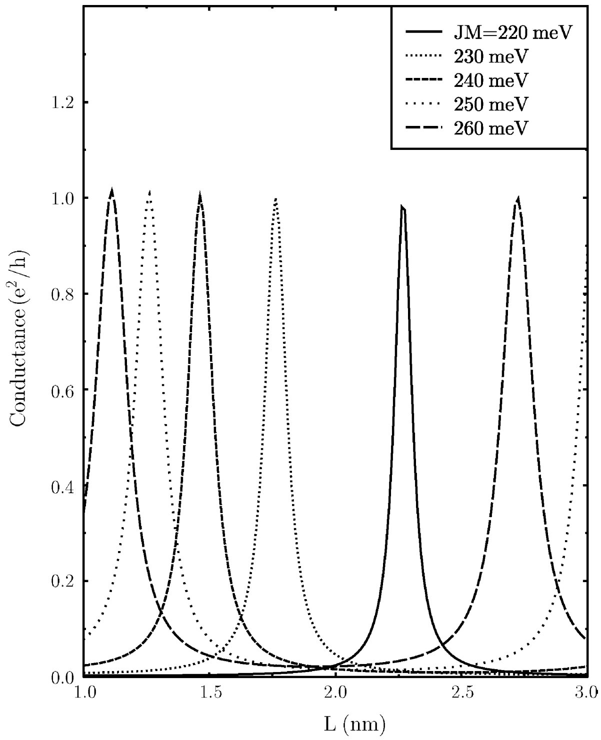

Fig. 2. Conductance vs. distance *L* between the two domain walls for different values of the magnetization strength *M*, but for a fixed carrier density.

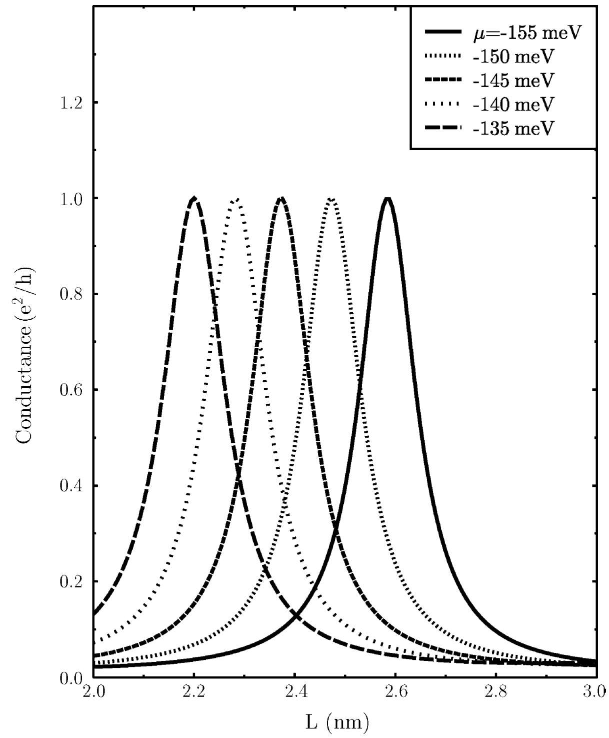

Fig. 3. Conductance vs. distance between two domain walls for different values of the chemical potential μ. The magnitude of magnetization *M* is constant.

coefficient t at the Fermi energy $\varepsilon = \mu$,

$$G = \frac{e^2}{2\pi\hbar}\,|t|^2. \tag{12}$$

Using the solution of Eqs. (7)–(10), we calculated the conductance as a function of the distance L for several values of the magnetization M. The results for $m = 0.6m_0$, $\delta = 2\times10^{-8}$ cm, and $\varepsilon = -200$ meV are presented in Fig. 2. for various values of M. As it is evident from Fig. 2, the conductance possesses narrow resonances which occur when the quasi-discrete energy levels in the spin quantum well coincide with the Fermi energy. In the vicinity of a resonance the conductivity changes drastically upon a small variation of the distance between the domain walls. Obviously, the same effect is achieved when, for a fixed L, the Fermi energy is changed, as demonstrated by Fig. 3, where we assumed $\delta = 2\times10^{-8}$ cm, and $JM = 220$ meV. Experimentally, one can achieve this by gating the whole structure, which leads to a change in the density of carriers and hence to a change in the position of the chemical potential. On the other hand, if one uses a tip or a gate to change the energy profile in the region $-L < z < L$ Fig. 1), one can then fix the location of the domain walls.

In this context we note that a similar effect of resonant tunnelling is realized in the resonance tunneling diode [15]. In such a device, the magnetic structure contains a thin layer, within which the quantizated energy levels are controlled by applying a magnetic field. However, in this case the width of the quantum well is not tunable.

4. Conclusions

We demonstrated theoretically the existence of sharp resonant peaks in the conductance of magnetic nanowires with a double domain walls. The peaks are controllable by changing characteristic parameters of the domain walls, such as the distance *L* between the walls or the density of carriers. The extreme sensitivity of the conductance to *L* can be used to identify the relative position of the domain walls. Alternatively, it can transform the effect of magnetic field on the domain wall into a current through the nanowire. We also demonstrated explicitly (Fig. 2) that a wall displacement of the order of 10% induces a change in the resistance up to hundreds of percents.

Acknowledgements

This work is supported by FCT Grant No. POCTI/FIS/58746/2004 (Portugal) and by Polish State Committee for Scientific Research under Grants Nos. PBZ/KBN/044/P03/2001 and 2 P03B 053 25. One of the authors (V.D.) thanks the Calouste Gulbenkian Foundation in Portugal for support.

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