Effect of bias on indirect exchange within magnetic nanostructures

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We formulate a theory of a reversible switching of the magnetic state of magnetic multilayers embedded in the metallic nanoconstriction by the external bias. The switching is related to the effect of strongly nonequilibrium electron distribution existing in the biased nanoconstriction on the indirect exchange coupling between the magnetic layers. © *2005 American Institute of Physics*. [DOI: 10.1063/1.2009810]

Switching devices based on the magnetic multilayers are recognized as one of the most promising for novel emerging computer and telecommunication technologies. Transport in these devices is controlled by coupling between magnetic layers, and the prospects for technological use are related to possibilities of easy tuning magnetic properties of multilayers. One of the key directions is the investigation of the particular heterostructures where, and means, by which the characteristics of the coupling itself can be controlled and tuned. The recent papers[1,2] proposed a novel concept of manipulating the magnetic configuration of heterostructures by an applied bias. They showed that in a magnetic bilayer separated by a thin insulating film, allowing tunneling current, the coupling via Ruderman–Kittel–Kasuya–Yosida (RKKY) interaction oscillates as a function of a bias. Thus, changing the applied voltage one can tune interaction going from ferromagnetic to antiferromagnetic coupling and vice versa.

However, a plateletlike geometry may not be the most beneficial for a practical design. Indeed, since tunneling current depends exponentially on the thickness of the insulating layer d, even the small variations in d over the area of the contact will destroy the coherent coupling. Maintaining layer thickness homogeneity of the order of the nearly atomic scale poses a serious challenge to technology. Additional restrictions to a useful range of parameters arise from the fact that coupling itself is exponentially small. Moreover, there is a controversy between the predictions by Refs. 2 and 1. While, according to Ref. 1, RKKY interaction in the structures in question does not depend on the thicknesse of the probe slab d and RKKI has a "surface" character as is the case of the standard equilibrium case,[2] finds the strength of RKKY in the nonequilibrium case to increase with increase of d (later, Ref. 3 arrived at the similar conclusions).

In this letter, we develop a theory of a tunable *point* spintronic device with the bias controllable exchange. This approach retaining a bias controlled exchange inherent to heterostructures, enables us at the same time to utilize the full power of quantum point contacts and thus allows us to avoid the above drawbacks of the plateletlike configuration. We consider a point contact device where nanometer-size ferromagnet plates are embedded into a narrow constriction between the two three-dimensional metallic electrodes (see Fig. 1). Such a configuration can support a highly nonequilibrium electronic distribution. Note that while the voltage drop is concentrated within the constriction region itself, the relaxation of the electron energy occurs over the distances well exceeding the constriction size. As a result, the point contacts can withstand fairly high biases, as compared to the Fermi energy, without disintegrating (see Ref. 4). We will also touch briefly on the contradictions between different papers addressing the switching of magnetic state of the hybrid ferromagnetic structures by the applied bias mentioned above.

We begin with a short diffusive channel of the length $\mathcal{L}$ enclosed between the two metal half-spaces implying that the electron mean-free path (assumed to be constant along the channel length) is much less than the channel length. It is known that in this case, the current-carrying part of the distribution function is much less than the part which is even in electron momentum. We also note that in this case the role of single acts of scattering, such as boundary scattering, can be neglected. Thus, the part even in the electron momentum can be found with a help of a diffusion equation $\Delta f=0$. The solution obeying the boundary conditions at the both ends of the channel and the neutrality condition $\int d\varepsilon f=\text{constant}$ can be written as

$$f=(1-x)F_0(\varepsilon_{\mathbf{k}}+eVx)+xF_0[\varepsilon_{\mathbf{k}}-eV(1-x)], \tag{1}$$

where F_0 is the Fermi function, $x=\tilde{x}/\mathcal{L}$, $\tilde{x}$ is the position within the channel, and V is the bias. Note that similar dis-

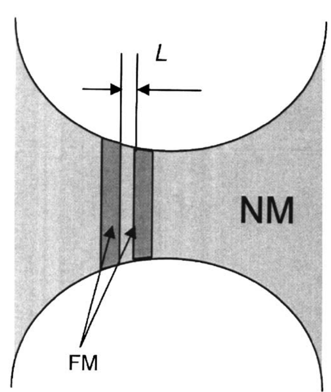

FIG. 1. Pointlike multilayer NM/FM/NM/FM/NM switching device.

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tribution function was exploited in Ref. 5. In what follows, we consider the configuration where the two thin ferromagnet layers, with thicknesses t_1 and t_2 and the spatial separation $L \ll \mathcal{L}$ are built in the channel made out of a nonferromagnetic metal. Here $(t_1, t_2) \ll \mathcal{L}$ and the layers do not significantly affect either the distribution of electric field or the electron distribution given by Eq. (1). The channel thickness exceeds the lattice constant a; this allows one to neglect the surface exchange energy and consider the magnetization within the layers as uniform.

The indirect exchange coupling between two ion spins has the form (see e.g., Ref. 6)

$$U_{\rm int} = 2j^2 \mathbf{S_i S_j} \sum_{\mathbf{k},\mathbf{k}'} \frac{(f_{\mathbf{k}} - f_{\mathbf{k}'}) \exp i(\mathbf{k}-\mathbf{k}')\mathbf{R_{ij}}}{\varepsilon_{\mathbf{k}} - \varepsilon_{\mathbf{k}'}}, \tag{2}$$

where j is an exchange energy. For the equilibrium at $T=0$, it gives

$$U_{\rm int} = \mathbf{S_i S_j} \frac{4j^2 m k_F^4 a^6}{(2\pi)^3 \hbar^2} F(2k_F R_{i,j}), \tag{3}$$

where

$$F(z) = \frac{\cos z}{z^3}.$$

Then, the coupling energy (per unit area) between the two ferromagnetic layers is obtained by the integration of Eq. (2) over the layers and reads[7]

$$\frac{E}{A} = \frac{j^2 m \mathbf{S_i S_j}}{16\pi^2 \hbar^2} K(k_F L), \tag{4}$$

where L is the spatial separation of the layers, and $K(z) \sim \sin z / z^2$. To take into account the effect of the nonequilibrium distribution given by Eq. (1) on the coupling, we rewrite relation (2) in terms of the scattering perturbation theory (see, e.g., Ref. 8) (that is using the scattering states as the first-order perturbative states):

$$U_{\rm int} = 2j^2 \mathbf{S_i S_j} a^3 \sum_{\mathbf{k}} \frac{1}{4\pi\varepsilon_k k} \frac{\exp i(\mathbf{k R_{ij}} - kR_{ij})}{R} f$$
$$= -2j^2 \mathbf{S_i S_j} \frac{a^3}{8\pi^3 \varepsilon_F k_F R^2} \int_0^\infty dk k f_k \sin(2kR_{ij}). \tag{5}$$

The indirect exchange interaction thus results from the Friedel oscillations of electronic density. Note that our approach discards the possibility of the net spin polarization of the electron current passing through the system when the voltage is applied. Below, we will briefly discuss the relation of our approach with other approaches exploiting the spin polarization in question either directly or indirectly.

Note that our calculations are restricted to the first Born approximation. Furthermore, the scheme discussed above discards the specifics of the ferromagnet wave functions assuming that the total phase difference is acquired within the nonmagnetic spacer. While both of the factors can affect, in principle, the phase of the Friedel oscillation, we do not expect it to affect the bias-dependent contribution into the phase since it is formed at large distances on the order of the thickness of the normal layer. We are interested in the effect of the bias which is not altered by the abovementioned simplifications.

After simple algebra, plugging in the finite temperatures expression (1) for f_k, one arrives at the following expression for the function $K(z)$:

$$K \sim \mathcal{G}\frac{1}{z^2}\left[\sin z\left(x\cos(1-x)z\frac{eV}{2\varepsilon_F} + (1-x)\cos xz\frac{eV}{2\varepsilon_F}\right)\right.$$
$$\left. + \cos z\left((1-x)\sin xz\frac{eV}{2\varepsilon_F} - x\sin(1-x)z\frac{eV}{2\varepsilon_F}\right)\right]. \tag{6}$$

The function $\mathcal{G}$ accounts for the finite temperatures:

$$\mathcal{G}(T) = \frac{4\pi m T L}{k_F \hbar^2} \sinh^{-1} \frac{4\pi L m T}{k_F \hbar^2}, \tag{7}$$

where $\mathcal{G}=1$ for $T \ll k_F \hbar^2 / 4\pi m L$ and exponentially decays for $T > k_F \hbar^2 / 4\pi m L$. At small V, the linear (in V) terms in Eqs. (5) and (6) vanish. Moreover, the expansion of the coefficient at the $\propto \cos z$ term in Eq. (6) starts from the third power. Thus, in the situation considered, bias has to be high enough $(eV/\varepsilon_F) k_F L \geq 1$ to cause a noticeable effect. This condition becomes less restrictive in ballistic systems where the distribution function $f(\mathbf{k})$ has a pronounced dependence not only of the magnitude of the wave vector, but also of its direction. In particular, in the ballistic limit, which is the most appropriate regime for the short channel or nanoconstriction, the distribution function for the central region has the form

$$f_{\mathbf{k}} = \theta(k_x) F_0\left(\varepsilon + \frac{eV}{2}\right) + \theta(-k_x) F_0\left(\varepsilon - \frac{eV}{2}\right), \tag{8}$$

where OX is the constriction axis. For the spherical Fermi surface at zero temperatures, this distribution corresponds to two semispheres: $k_x > 0$ and $k_x < 0$, with radii differed by $\Delta k = k_F (eV/2\varepsilon_F)^{1/2}$. Note that such a solution was first found in Ref. 9. Let us assume that two of the ferromagnetic layers are placed in the central region in question while the thickness of the region occupied by the layers is small enough to neglect the spatial dependence of the distribution function. We also assume that the interfaces are perfect and do not affect the distribution of Eq. (8) significantly. Indeed, actually it holds at least for the modes corresponding to trajectories passing through the contact without backscattering.

In what follows, we will restrict ourselves to the case of small biases $eV < \varepsilon_F / k_F R$. Then one can linearize the expression in Eq. (8) and get

$$f_{\mathbf{k}} = F_0 + \frac{eV}{2} \frac{\partial F_0}{\partial \varepsilon} \mathrm{sign}(k_x). \tag{9}$$

The integration over $\mathbf{k}$ in Eq. (5) is reduced then to

$$-\mathcal{G}\frac{k_F}{2R^2}\cos(2k_F R)$$
$$+ \frac{eV}{2} \sum_{\mathbf{k}_\perp, k_x > 0} \frac{\sin k_x R_x \sin(\mathbf{k}_\perp \mathbf{R}_\perp - kR)}{4T \cosh^2[(k_F(k-k_F)\hbar^2/4mT)]}. \tag{10}$$

Here the subscript $\perp$ denotes components normal to the contact axis. It is important to bear in mind that in the situation considered the effect of the electron system on the ion spins cannot be reduced to the spin coupling. Indeed, the result of the integration in Eq. (10) will depend on the direction of the vector $\mathbf{R}$ connecting ith and jth spins. In other words, the effect of spin i on the spin j is different from the effect of spin j on spin i. It is not surprising since we deal with a strongly nonequilibrium electron distribution and the effect of this distribution on the spins cannot be described as a simple screening. Rather, it is a transfer of spin polarization by the current flux. To simplify the picture, we will assume that the magnetization direction for one of the ferromagnets is fixed and that we deal therefore with the bias-induced transfer of the polarization to the second ferromagnet.

As a result, one finally obtains

$$K \sim \frac{1}{z^2}\left(\mathcal{G}(T)\sin z - z\frac{eV}{2\varepsilon_F}\mathcal{G}(T/2)\mathcal{K}_1(z)\right), \qquad (11)$$

where $\mathcal{K}_1(z)$ is the rapidly oscillating function of the order of unity. According to our results, the application of the bias can change the mutual orientation of magnetizations of the two ferromagnetic layers coupled within the constriction. Let the easy magnetization axis of the two ferromagnets be parallel (this occurs if $k_F L = \pi n$). In this case, in the absence of coupling between the layers and/or if external magnetic field is zero, the parallel or antiparallel alignment of the magnetizations are equally probable. The external bias lifts this symmetry, thus making one of the configurations more favorable, depending on the sign of the bias. Thus one can, in a controllable way, switch between the two configurations applying the external voltage. While the pulses of large bias can operate the system ("writing"), the small current through the system can be used for reading with the help of the giant magnetoresistance effect, making the resistance depend upon the mutual orientation of the magnetization vectors.

Now, let us compare the present mechanism of switching with those implying a presence of the net spin polarization of the current through the structure. Among them, one notes the papers[1,3] discussing "nonequilibrium exchange interaction" and the papers[10,11] considering the spin transfer from the incident spin current to the ferromagnet. The apparent dependence of the nonequilibrium RKKY on t_2 stated in Refs. 1 and 3 is quite disputable. Namely, according to Ref. 3, the t_2 dependence eventually follows from the fact that the total spin of the slab is coupled to a net spin transported by the current through the structure. The essential assumption of Ref. 3 that this net spin is completely controlled by the polarizing slab and is not affected by the probe slab, except that spin relaxation is not justified. Indeed, coupling of the spin of an electron passing through the probe slab to the magnetization within the slab leads to a precession of the electron spin. Moreover, due to the electron motion through the probe slab the spatial evolution of such a precession is different for electrons with different momenta. As a result, the spin current effectively affects the total spin of the slab only at distances smaller than spin precession length $l_{ex} \sim (\hbar/p_F) \times(\varepsilon_F/E_{ex})$, where E_{ex} is the exchange energy. Therefore, the effect of the spin current depends upon the slab thickness only if $t_2 < l_{ex}$. It seems that the contribution discussed in Refs. 1 and 3 is of the same nature as that considered by Ref. 10 and effectively reduces to a surface "torque." Consequently, this contribution is not of the same nature as true RKKY interactions related to spatial oscillations of the electron density.

References 10 and 11 directly exploit the spin pumping from one ferromagnet to another. While the mere existence of such a mechanism may be viewed as well established, its efficiency is to be thoroughly examined since it depends critically on the ability of one of the ferromagnets involved to polarize the current passing through the system. On the contrary, the mechanism proposed here does not require spin polarization of the current and, therefore, is expected to dominate the behavior of the thin ferromagnetic layers devices where both t_1 and t_2 are small. Another note is that although in the above discussion we focused on the setup where ferromagnets are normal to the constriction axis, the arrangement when they are parallel to the constriction axis is also possible since our mechanism is based on the nonequilibrium distribution function rather than on the current through the ferromagnets. In this latter, case no spin polarization can occur.

To summarize, we have shown that magnetic state of magnetic heterostructures embedded into metallic nanoconstrictions can be reversibly switched by a pulses of the applied bias of different signs. The switching takes place due to an effect of strongly nonequilibrium electron distribution on the indirect exchange between the magnetic layers which is related to a bias dependence of the Fermi momentum.

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