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Resonant magnetic tunnel junction at 0° K: I-V characteristics and magnetoresistance

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In this paper we analyze the main transport properties of a simple resonant magnetic tunnel junction (FM-IS-METAL-IS-FM structure) taking into account both elastic and magnon-assisted tunneling processes at low voltages and temperatures near 0° K. We show the possibility of magnetoresistance inversion as a consequence of inelastic processes and spin-dependent transmission coefficients. Resonant tunneling can also explain the effect of scattering by impurities located inside an insulating barrier.
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I. INTRODUCTION

Recently, the interest on the phenomena of Giant Magnetoresistance (GMR) in magnetic tunnel junctions (MTJ), made up of two ferromagnetic electrodes separated by an insulating barrier, has grown significantly due to potential applications in magnetoresistive reading heads, magnetic field sensors, nonvolatile magnetic random access memories, and many others.[1–5] The GMR effect is based on the spin-dependent scattering mechanisms proposed in the early papers by Cabrera and Falicov,[6] which lead in MTJ's, to a strong dependence of the conductance on the magnetic polarization.[7] Typically, the GMR effect found in MTJ's is of the order of 25%–30%,[8–10] and points to a large ratio of the densities of states for majority (M) and minority (m) electrons at the Fermi level (E_F),[8,11,12]

$$\frac{N_M(E_F)}{N_m(E_F)} \approx 2.0-2.5.$$

As usual in MR experiments, one compares the resistances for the cases where the magnetizations at the electrodes are antiparallel (AP) and parallel (P).[13]

The subject of this paper is a quite similar ferromagnetic structure, the resonant magnetic tunnel junction (RMTJ), made up of two ferromagnetic electrodes with a layered insulator–metal–insulator structure between them. The resonance energy levels also can be provided by impurity scattering centers placed inside a single insulating barrier.[14–16] Experimentally, it was found that impurities in the barrier[17] and electronic structure of the interfaces strongly affect the MR.[18] A remarkable effect of resonant tunneling in magnetic structures is the inversion of spin-polarization,[19] with a consequent inversion of the MR.

In Ref. 8, scattering from magnons at the electrode–insulator interface has been proposed as the mechanism for randomizing the tunneling process and opening the spin-flip channels that reduce the MR in single barrier MTJs. This process may explain the MR behavior in the vicinity of zero bias (voltages smaller than 40–100 mV). In fact, inelastic-electron tunneling spectroscopy (IETS) measurements at low temperature showed peaks that can unambiguously be associated with one-magnon spectra at very small voltages (from 12 to 20 mV, with tails up to 40 mV, and maximum magnon energy not larger than 100 meV[1]). In the present work the resonant MTJ constituted of a double barrier structure will be analyzed, and the inelastic magnon scattering will be proposed as a mechanism capable to produce the inversion of MR. A complete theory of the transport properties of a RMTJ should include the following:

(i) magnon assisted tunneling effects, with maximum magnon energies of the order of ~100 meV. At low temperature, electrons from the electrodes, accelerated by the applied voltage, excite magnons at the interface. At low temperature, only magnon-emission processes should be considered;

(ii) variation with voltages of the densities of states for the different spin bands in the ferromagnets. Here, we will follow closely the approach of Refs. 11 and 12, with a simple picture of the band structure. This is motivated by the discussions given in Ref. 11 over the polarization of the tunneling current. We assume here that the latter is mainly of s character;

(iii) consideration of a resonant tunneling transmission coefficient with *Lorentzian* shape near the resonance, and voltage dependence of the resonant energy level. The resonances for elastic and magnon-assisted processes are different, allowing the inversion of MR.

The above program will be developed in the present contribution. The content of this paper can be described as follows: In the next section, we formulate the theoretical basis for analyzing tunneling currents, discussing the transfer Hamiltonian that includes all the above-mentioned ingredients. In Sec. III, we analyze the I-V characteristics and the MR inversion considering resonant transmission coefficients and the effects of density of states in both elastic and magnon-assisted tunneling processes. Finally, in the last section, a few conclusions and remarks are added.

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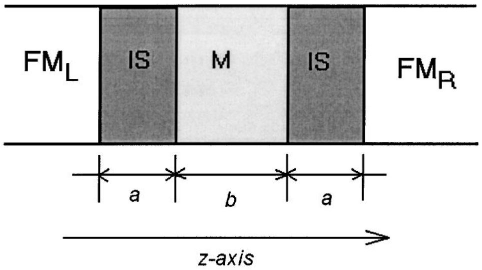

FIG. 1. A simple double barrier resonant magnetic tunnel junction: FM_I and FM_{II} are two ferromagnetic electrodes, IS_I and IS_{II} are two potential barriers while M is a metallic layer (ferromagnetic or not) placed between the two barriers.

II. THEORETICAL FRAMEWORK

To give a description of the MR and the resistance in a RMTJ, we will use the transfer Hamiltonian method.[20–22] The resonant junction is shown schematically in Fig. 1 FM_L and FM_R are two ferromagnetic electrodes. The regions designated by IS are two insulators (thin oxide films) representing a potential barrier for conduction electrons due to the fact that the Fermi levels of the ferromagnetic layers are situated in the gap region of the oxide films. The region designated by M is a metallic layer (which can be ferromagnetic or not) located between the two insulating regions and having the Fermi level similar to that of the ferromagnetic electrodes FM_L and FM_R. We have considered the s-band electrons as free particles (plane waves), being responsible for the dominant contribution to the tunneling process. The d electrons, which are more localized, enter in the process only via the exchange interaction with s electrons on each ferromagnetic electrode. In the context of second quantization and neglecting the magnetization energy (Zeeman term), the unperturbed Hamiltonian, representing the ferromagnetic reservoirs FM_L and FM_R, is given by

$$H_0 = \sum_{\mathbf{k}\sigma,\alpha=(L,R)} E_{\mathbf{k}\sigma} c^{\alpha\dagger}_{\mathbf{k}\sigma} c^{\alpha}_{\mathbf{k}\sigma}, \tag{1}$$

with $L(R)$ referring to the left(right) ferromagnetic electrode, $c^{\alpha\dagger}_{\mathbf{k}\sigma}$ $(c^{\alpha}_{\mathbf{k}\sigma})$ are the creation (annihilation) fermionic operators for wave vector $\mathbf{k}$ and spin σ, $E_{\mathbf{k}\sigma}=\hbar^2k^2/2m-\sigma\Delta_\alpha$ is the Hartree–Fock energy, and Δ_α is the shift in energy due to the exchange interaction on each side of the double barrier.

The double barrier structure (IS-M-IS) can be treated as a perturbation, being responsible for the interaction between the two ferromagnetic electrodes. In writing the interaction part of the total Hamiltonian, which makes possible the transfer of electrons from one side ferromagnetic electrode to the other, we follow Ref. 8. Apart from the direct transfer that comes from elastic processes, we include transfer with magnetic excitations that originates from the s-d exchange between conduction electrons and localized spins at the two interfaces FM_L-IS and IS-FM_R. The excitations are described by a linearized Holstein–Primakoff transformation,[22] in the spirit of a one-magnon theory. We use the following interaction Hamiltonian:

$$\begin{aligned} H_I = {} & \sum_{\mathbf{kk}'\sigma} t^d_{\mathbf{kk}'} (c^{L\dagger}_{\mathbf{k}\sigma} c^R_{\mathbf{k}'\sigma} + c^{R\dagger}_{\mathbf{k}'\sigma} c^L_{\mathbf{k}\sigma}) + \frac{1}{\sqrt{N_s}} \sum_{\mathbf{kk}'\mathbf{q}} t^J_{\mathbf{kk}'\mathbf{q}} (c^{L\dagger}_{\mathbf{k}\downarrow} c^R_{\mathbf{k}'\uparrow} \\ & + c^{R\dagger}_{\mathbf{k}'\downarrow} c^L_{\mathbf{k}\uparrow})(\sqrt{2S_L} b^L_{\mathbf{q}} + \sqrt{2S_R} b^R_{\mathbf{q}}) \\ & + \frac{1}{\sqrt{N_s}} \sum_{\mathbf{kk}'\mathbf{q}} t^J_{\mathbf{kk}'\mathbf{q}} (c^{L\dagger}_{\mathbf{k}\uparrow} c^R_{\mathbf{k}'\downarrow} + c^{R\dagger}_{\mathbf{k}'\uparrow} c^L_{\mathbf{k}\downarrow})(\sqrt{2S_L} b^{L\dagger}_{\mathbf{q}} \\ & + \sqrt{2S_R} b^{R\dagger}_{\mathbf{q}}) + \frac{1}{N_s} \sum_{\mathbf{kk}'\mathbf{q}} t^J_{\mathbf{kk}'\mathbf{q}} (c^{L\dagger}_{\mathbf{k}\uparrow} c^R_{\mathbf{k}'\uparrow} - c^{L\dagger}_{\mathbf{k}\downarrow} c^R_{\mathbf{k}'\downarrow} + \text{h.c.}) \\ & \times (S_L + S_R - (b^{R\dagger}_{\mathbf{q}} b^R_{\mathbf{q}} + b^{L\dagger}_{\mathbf{q}} b^L_{\mathbf{q}})), \end{aligned} \tag{2}$$

where $t^d_{\mathbf{kk}'}$ is the direct transmission amplitude, $t^J_{\mathbf{kk}'\mathbf{q}}$ is the inelastic transmission amplitude (depends on the exchange integral), $S^L(S^R)$ is the spin value at the left (right) ferromagnetic electrode, N_s is the total number of spins at the interfaces FM_L-IS and IS-FM_R, and $b^{\alpha\dagger}_{\mathbf{q}}$ $(b^{\alpha}_{\mathbf{q}})$ are the creation (annihilation) operators for magnons with wave vector $\mathbf{q}$ at each interface between the double barrier IS-M-IS and the electrodes FM_L and FM_R. The wave vector $\mathbf{q}$ is quasi-two-dimensional (the magnon wave function is localized at the interfaces, but with finite localization length). The origin of the direct and inelastic transmission amplitudes, $t^d_{\mathbf{kk}'}$ and $t^J_{\mathbf{kk}'\mathbf{q}}$, respectively, is the overlap between the left and right electrode wave functions inside the double barrier region, defined simply as follows:

$$t^d_{\mathbf{kk}'} = \int d^3\mathbf{x}\ \varphi^{L*}_{\mathbf{k}}(\mathbf{x}) \left(\frac{\mathbf{p}^2}{2m} + V(\mathbf{x}) \right) \varphi^R_{\mathbf{k}'}(\mathbf{x})$$

and

$$t^J_{\mathbf{kk}'\mathbf{q}} = -\int d^3\mathbf{x} \int d^3\mathbf{x}'\ J(\mathbf{x},\mathbf{x}')\varphi^{L*}_{\mathbf{k}}(\mathbf{x})\varphi^R_{\mathbf{k}'}(\mathbf{x})\exp[i\mathbf{q}\cdot\mathbf{x}'],$$

$\varphi^{\alpha}_{\mathbf{k}}(\mathbf{x})$ being the α-electrode wave function, $V(\mathbf{x})$ the applied voltage superimposed to the lattice potential, and $J(\mathbf{x},\mathbf{x}')$ the position-dependent exchange factor. The transmission coefficients will be proportional to the above quantities, $T^d_{\mathbf{kk}'} = |t^d_{\mathbf{kk}'}|^2$ and $T^J_{\mathbf{kk}'\mathbf{q}} = |t^J_{\mathbf{kk}'\mathbf{q}}|^2$. As one can clearly see from (2), the direct tunneling is responsible for a conserving spin current while the inelastic tunneling provides a spin-flipping magnon-assisted current flowing through the RMTJ. A rigorous analysis of the transmission coefficients for a double barrier structure placed between two metallic ferromagnetic electrodes is given in Ref. 20, showing that the transmission coefficients, $T^d=|t^d|^2$ and $T^J=|t^J|^2$, close to the resonance levels, are very well fitted through the use of a *Lorentzian* function, defined as follows:

$$T(\epsilon) = \frac{\Gamma^2(\epsilon)}{\Gamma^2(\epsilon) + (\epsilon - \epsilon_r(V))^2}, \tag{3}$$

where

$$\Gamma(\epsilon) = \left[\frac{\hbar^2(\epsilon_0 + |\epsilon_r(V)|)}{4meb^2} T_1^2(\epsilon) \right]^{1/2} \tag{4}$$

is the half-width at the half-maximum of the resonance in the potential well formed by the structure IS-M-IS, m is the electron mass, b is the metallic layer thickness, $T_1(\epsilon)$ is the transmission coefficient for a single barrier, which is given below, in the WKB approximation:

$$T_1(\epsilon) \approx A \exp[-2a\sqrt{2me(V_0-\eta\epsilon)/\hbar^2}] = A \exp[-1.024d\sqrt{V_0}]\exp\left[\frac{1}{2}\frac{\eta\epsilon a}{\sqrt{V_0}}\right],$$

with A constant, $\eta=\epsilon_z/\epsilon$, relating the energy ϵ with its component ϵ_z perpendicular to the double barrier, a is the insulating barrier thickness in Angstroms, ϵ_0 and $\epsilon_r(V)$ are related to the resonant energy, $\epsilon_r(V)$ being a function of the applied voltage V,

$$\epsilon_r = \epsilon_0 - \frac{V}{2} \quad (5)$$

and all energies given in eV. The above formulas, shown in energy variables, for the sake of convenience, will be used later on in this paper.

In general, the total current obtained with (2) has contributions from elastic processes, resulting in a direct tunneling that conserves spin, and from the inelastic ones, which involve the emission and absorption of magnons with electronic spin flip. In the following we describe the transport properties of the RMTJ.

III. TRANSPORT PROPERTIES: I-V CHARACTERISTICS AND MR INVERSION

The interaction Hamiltonian (2) and the Fermi Golden Rule,

$$w_{\mathbf{k}\sigma-\mathbf{k}'\sigma'} = \frac{2\pi e}{\hbar}|\langle\mathbf{k}'\sigma'|H_I|\mathbf{k}\sigma\rangle|^2\delta(E_{k\sigma}-E_{k'\sigma'}\pm\hbar\omega), \quad (6)$$

allow us to express the tunneling current as follows:

$$I = \frac{2\pi e}{\hbar}\sum_{\mathbf{k}\mathbf{k}'\sigma\sigma'}|\langle\mathbf{k}'\sigma'|H_I|\mathbf{k}\sigma\rangle|^2 \times (f(E^L_{\mathbf{k}\sigma})[1-f(E^R_{\mathbf{k}'\sigma'})] - f(E^R_{\mathbf{k}'\sigma'})[1-f(E^L_{\mathbf{k}\sigma})])\delta(E_L-E_R). \quad (7)$$

Now we replace the sums over the wave vectors $\mathbf{k}$ and $\mathbf{k}'$ by energy integrals, which is a well-known procedure, integrate the Dirac delta functions $(\delta(E_L-E_R))$, and obtain the total current flowing through the junction

$$I(V) = \frac{2\pi e}{\hbar}\int d\epsilon\sum_{\sigma\sigma'}(T_{\sigma\sigma'}(\epsilon)N^L_\sigma(\epsilon-V)N^R_{\sigma'}(\epsilon)f_L(\epsilon-V) \times[1-f_R(\epsilon)] - T_{\sigma'\sigma}(\epsilon)N^L_\sigma(\epsilon-V)N^R_{\sigma'}(\epsilon)f_R(\epsilon) \times[1-f_L(\epsilon-V)]), \quad (8)$$

being $\epsilon=E-E_F$, E_F the Fermi energy, $N^\alpha_\sigma(\epsilon)$ the density of states for electrons with spin σ at the electrode FM_α, $\alpha=(R,L)$ denote the electrode, and $f(\epsilon)=(\exp[\epsilon/k_BT]+1)^{-1}$ is the Fermi–Dirac distribution, k_B is the Boltzmann constant, and T the absolute temperature. The tunneling coefficients are obtained from (2) through $T_{\sigma\sigma'}(\epsilon)=|\langle\mathbf{k}\sigma|H_I|\mathbf{k}'\sigma'\rangle|^2$, being related to the direct and inelastic transmission coefficients by $T_{\sigma\sigma}=T^d$ and $T_{\sigma,-\sigma}=T^J$, respectively. These quantities are functions of energy, applied voltage, and double barrier parameters.

We are interested now in the basic transport properties such as the I-V characteristic, the differential conductance $G=dI/dV$ and the MR. The tunneling magnetoresistance (TMR) is defined as

$$\mathrm{TMR} = \frac{I^{(P)}-I^{(AP)}}{I^{(P)}}, \quad (9)$$

where P(AP) stand for the parallel(antiparallel) configuration scheme. In the P configuration the magnetization (imposed by external applied magnetic fields) of both electrodes FM_L and FM_R has the same orientation along the z axis, the majority and minority spin bands corresponding to the same spin orientation for both electrodes, while in the AP configuration the magnetization is rotated by 180° from one electrode to another, and in this case, the majority spin band in FM_L is the minority spin band in FM_R and *vice versa*. In the low bias regime, we are interested in voltages smaller than the Fermi energy and only the states near the Fermi level will contribute to the transport, so we can expand the density of states in a Taylor series as follows:

$$N^\alpha_\sigma(\epsilon) = \sum_{n=0}^{\infty}\frac{1}{n!}\left.\frac{d^nN^\alpha_\sigma(\epsilon)}{d\epsilon^n}\right|_{\epsilon=0}\epsilon^n. \quad (10)$$

Taking into account identical electrodes and the low bias regime, we can expand these expressions to first order with good accuracy. The s band can be represented by a parabolic dispersion relation and density of states $N_\sigma \propto \sqrt{E-\Delta_\sigma}$, where $\Delta_\sigma(\sigma=\uparrow,\downarrow)$ gives the bottom of the spin band, with $|\Delta_\uparrow-\Delta_\downarrow|=2\Delta$, as in Ref. 15. However, we consider here cases more general than the parabolic dispersion, with the band structure described through the following set of parameters:

$$r \equiv \left(\frac{N_M}{N_m}\right)_F, \quad \lambda \equiv \left(\frac{dN_M/dE}{dN_m/dE}\right)_F, \quad \beta \equiv \left(\frac{1}{N_m}\frac{dN_m}{dE}\right)_F, \quad (11)$$

with all quantities evaluated at the Fermi level. The labels m and M stand for minority and majority spin bands, respectively. It was shown in a previous work[16] that the ratio of the densities of states r at the Fermi level can be calculated using the experimental value of $\Delta R/R$ at zero bias (only if $T_{\sigma\sigma'}$ is almost independent of σ and σ' close to the voltage $V=0$):

$$r = \frac{1}{1-\left.\frac{\Delta R}{R}\right|_{V=0}} + \sqrt{\frac{1}{\left(1-\left.\frac{\Delta R}{R}\right|_{V=0}\right)^2}-1}, \quad (12)$$

which does not depend on the barrier parameters.

The total current given by (8) can be decomposed into a conserving spin direct current and a spin-flipping inelastic current involving magnon scattering. Considering symmetric transmission coefficients (equal probabilities for tunneling from the left to the right electrode as in the opposite sense), the direct current is written below:

$$I_d = \frac{2\pi e}{\hbar}\int_0^V d\epsilon (T_{\uparrow\uparrow}(\epsilon)N_\uparrow^L(\epsilon - V)N_\uparrow^R(\epsilon) + T_{\downarrow\downarrow}(\epsilon)N_\downarrow^L(\epsilon - V)N_\downarrow^R(\epsilon)). \quad (13)$$

The transmission coefficients ($T^d_{\sigma\sigma}$) are modeled according to the expression (3), as follows:

$$T_{\uparrow\uparrow} = \frac{\Gamma^2_{\uparrow\uparrow}(\epsilon)}{\Gamma^2_{\downarrow\downarrow}(\epsilon) + (\epsilon - \epsilon_{r\uparrow\uparrow}(V))^2},$$

$$T_{\downarrow\downarrow} = \frac{\Gamma^2_{\uparrow\uparrow}(\epsilon)}{\Gamma^2_{\downarrow\downarrow}(\epsilon) + (\epsilon - \epsilon_{r\downarrow\downarrow}(V))^2},$$

with $\epsilon_{r\uparrow\uparrow} \neq \epsilon_{r\uparrow\uparrow}$ in general. To evaluate the expression (13) we define the following quantities, depending only on density of states:

$$W_1(V) = N_M^L(\epsilon - V)N_M^R(\epsilon)|_{\epsilon=V/2} = (N_m^F)^2\left(r^2 - \frac{\beta^2\lambda^2}{4}V^2\right), \quad (14)$$

$$W_2(V) = N_m^L(\epsilon - V)N_m^R(\epsilon)|_{\epsilon=V/2} = (N_m^F)^2\left(1 - \frac{\beta^2}{4}V^2\right), \quad (15)$$

$$W_3(V) = N_M^L(\epsilon - V)N_m^R(\epsilon)|_{\epsilon=V/2} = (N_m^F)^2\left(r + \frac{\beta(r-\lambda)}{2}V - \frac{\beta^2\lambda}{4}V^2\right), \quad (16)$$

$$W_4(V) = N_m^L(\epsilon - V)N_M^R(\epsilon)|_{\epsilon=V/2} = (N_m^F)^2\left(r - \frac{\beta(r-\lambda)}{2}V - \frac{\beta^2\lambda}{4}V^2\right), \quad (17)$$

and the quantity $\Sigma_{\sigma\sigma'}(V)$, which is only dependent on the transmission coefficients, as follows:

$$\Sigma_{\sigma\sigma'}(V) = \int_0^V d\epsilon\, T_{\sigma\sigma'}(\epsilon) = \int_0^V d\epsilon \frac{\Gamma^2_{\sigma\sigma'}(\epsilon)}{\Gamma^2_{\sigma\sigma'}(\epsilon) + (\epsilon - \epsilon_{r\sigma\sigma'}(V))^2} = \Gamma_{\sigma\sigma'}(V)\left[\arctan\left(\frac{V - \epsilon_{r\sigma\sigma'}(V)}{\Gamma_{\sigma\sigma'}(V)}\right) + \arctan\left(\frac{\epsilon_{r\sigma\sigma'}(V)}{\Gamma_{\sigma\sigma'}(V)}\right)\right]. \quad (18)$$

Making use of (14)–(18), we obtain the explicit expressions for the conserving spin direct current, in both P and AP configurations:

$$I_d^{(\mathrm{P})} = \frac{2\pi e}{\hbar}[W_1(V)\Sigma_{\uparrow\uparrow}(V) + W_2(V)\Sigma_{\downarrow\downarrow}(V)], \quad (19)$$

$$I_d^{(\mathrm{AP})} = \frac{2\pi e}{\hbar}[W_3(V)\Sigma_{\uparrow\uparrow}(V) + W_4(V)\Sigma_{\downarrow\downarrow}(V)]. \quad (20)$$

When the spin flipping magnon-assisted current can be neglected, the MR inversion (the MR reverses its sign) is allowed in an interval of voltages if the requirement $I_d^{(\mathrm{P})}(V) < I_d^{(\mathrm{AP})}(V)$ is satisfied, i.e.,

$$\frac{\Sigma_{\downarrow\downarrow}(V)}{\Sigma_{\uparrow\uparrow}(V)} > \frac{W_1(V) - W_3(V)}{W_4(V) - W_2(V)}. \quad (21)$$

The above case is illustrated in Fig. 2, where we have neglected the inelastic current. We have used the following parameters: $r=2.21$ (from Refs. 8, 11, and 12), $\lambda=2.9$, $\beta=0.08$, $a=0.5$ nm, $b=0.1$ nm, $V_0=3.0$ eV. For Figs. 2(a)–2(c), we have chosen the resonant energies to be $\epsilon_{0\uparrow\uparrow}=35$ meV and $\epsilon_{0\downarrow\downarrow}=47$ meV, for spin-up and spin-down, respectively, while for Figs. 2(d)–2(f) we invert the situation making $\epsilon_{0\uparrow\uparrow}=47$ meV and $\epsilon_{0\downarrow\downarrow}=35$ meV. Is is clear from Figs. 2(a) and 2(d) that for a finite interval of voltages the current I_{AP} is greater than I_{P} with the consequent inversion of MR. Looking at Figs. 2(b) and 2(e), we observe that the differential conductance $G=dI/dV$ assume negative values for a range of voltages where the increase of applied voltage produce a reduction of the total current. This behavior is a consequence of the resonant tunneling known as Negative Differential Resistance (NDR) phenomena and takes place in several devices such as tunnel diodes.[20] Our attention turns now to Figs. 2(c) and 2(f), where we have plotted the TMR using the expression (9). Observe that in the case of Fig. 2(c) the MR starts from its zero bias value (25%) and increase with the applied voltage until reaching its maximum value near 50%. From its maximum value the MR begins to decrease with the increase of the applied voltage, assuming negative values and reaching a minimum. After that the MR turns to increase with the applied voltage to assume positive values again. Two important consequences follow from the resonant tunneling: (1) greater values of MR are obtained using the RMTJ (maximum MR about 50% against the 25% obtained in a simple MTJ; and (2) inversion of MR in a finite range of the applied voltage. Now we look at the situation shown in Fig. 2(d). The MR starts from a negative value (about −5%) near $V=0$, decreasing with the applied voltage until reaching a minimum value (∼−50%) to begin to grow up until becoming positive. The maximum positive value is of the order of 50% and the main consequences (an increase of MR values when compared to the simple MTJ's and the MR inversion) are achieved as in the previous case. The MR behavior of both situations can be easily understood with a basis in the resonant transmission coefficients and in the band structure. In the first situation [Figs. 2(a)–2(c)] the resonance for the spin-up transmission coefficient occurs first before for spin-down. When the spin-up resonance is the dominant contribution $[\Sigma_{\uparrow\uparrow}(V) > \Sigma_{\uparrow\uparrow}(V)]$ to the direct currents, the parallel I_{P} depends mainly on the product of majority spin bands $W_1(V) \propto N_M N_M$ while the antiparallel I_{AP} depends on $W_3(V) \propto N_M N_m$, being $W_1 > W_3$. The MR increase with the applied voltage until the spin-down resonance is reached while the spin-up resonance is turned off. Now that the spin-down channel is the dominant contribution to the

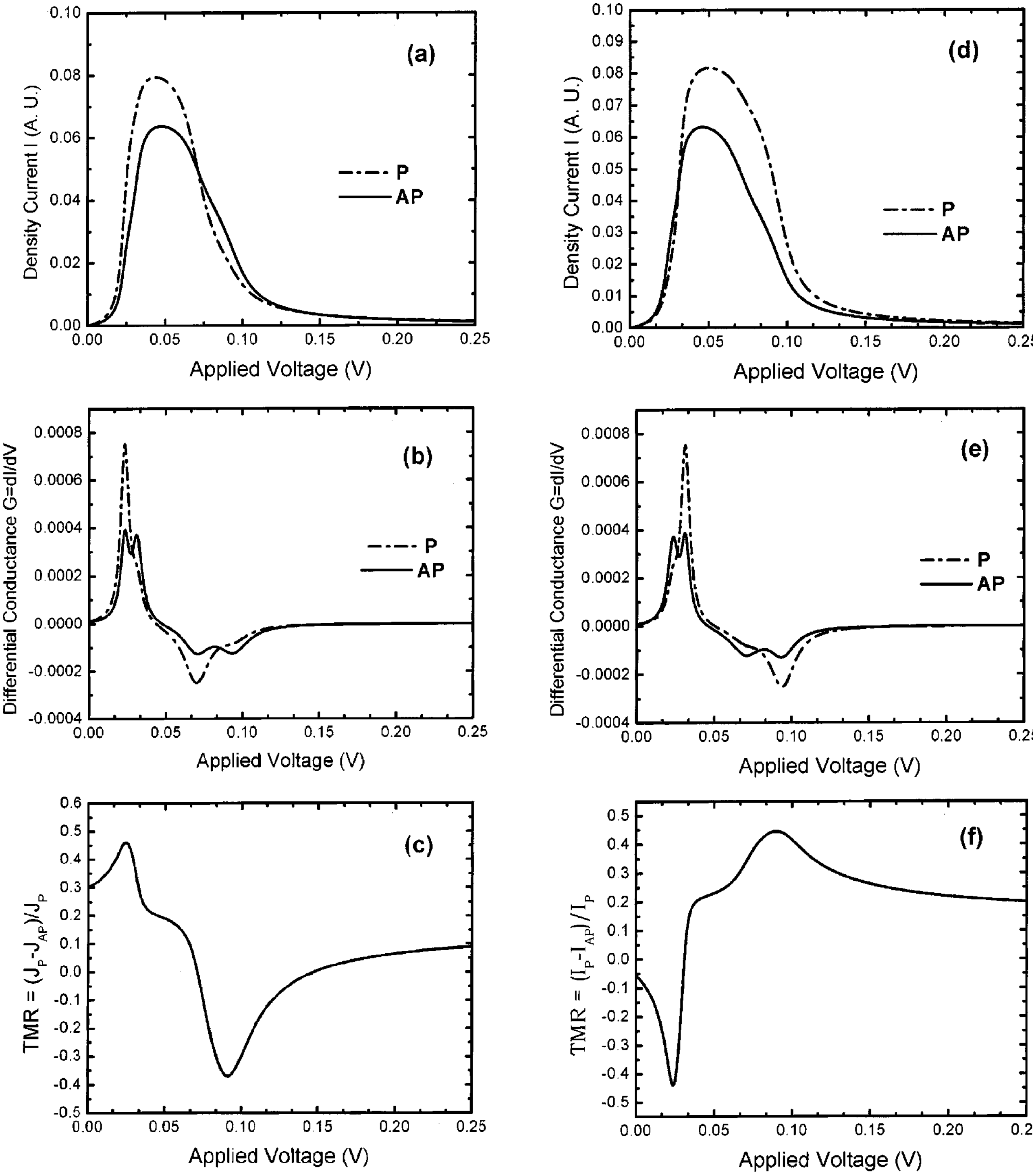

FIG. 2. Neglecting magnon-assisted tunneling we show (a) I-V characteristic; (b) differential conductance $G=dI/dV$ for both P and AP alignments; and (c) the tunneling magnetoresistance defined as $\mathrm{TMR}=(I_{\mathrm{P}}-I_{\mathrm{AP}})/I_{\mathrm{P}}$ using $\epsilon_{0\uparrow}=35$ meV and $\epsilon_{0\downarrow}=47$ meV (d), (e), and (f) show the same as (a), (b), and (c), respectively, for $\epsilon_{0\uparrow}=47$ meV and $\epsilon_{0\downarrow}=35$ meV, keeping the other parameters the same.

currents, the parallel I_{P} depends mainly on the product of minority spin bands $W_2(V)\propto N_M N_M$ while the antiparallel I_{AP} depends on $W_4(V)\propto N_m N_M$, being $W_4 > W_2$. In this way $I_{\mathrm{AP}} > I_{\mathrm{P}}$ and the MR inversion is achieved. For greater voltages both resonances are turned off and the resonant junction behaves as a simple MTJ. In the second situation, corresponding to Figs. 2(a)–2(c), spin-down resonance is reached first before spin-up resonance and the behavior is obvious.

Consider now that spin-up and spin-down resonances occur at the same energy, $\epsilon_{0\uparrow\uparrow}=\epsilon_{0\downarrow\downarrow}=\epsilon_{0d}$, which means $\Sigma_{\downarrow\downarrow}(V)=\Sigma_{\uparrow\uparrow}(V)=\Sigma(V)$, explicitly written below:

$$\Sigma(V) = \int_0^V d\epsilon T_{\uparrow\uparrow}(\epsilon) = \int_0^V d\epsilon T_{\downarrow\downarrow}(\epsilon) = \Gamma_d(V) \times\left[\arctan\left(\frac{V-\epsilon_{rd}(V)}{\Gamma_d(V)}\right) + \arctan\left(\frac{\epsilon_{rd}(V)}{\Gamma_d(V)}\right)\right]. \quad (22)$$

The band structure parameters $W_i(V)$ are well-behaved functions of the applied voltage V being the MR inversion forbidden through elastic processes (unless that $r<1$, which is impossible; by definition; r is always greater than unity). The direct currents in both P and AP configurations are given by

$$I_d^{(\mathrm{P})} = \frac{2\pi e}{\hbar}[W_1(V) + W_2(V)]\Sigma(V), \tag{23}$$

$$I_d^{(\mathrm{AP})} = \frac{2\pi e}{\hbar}[W_3(V) + W_4(V)]\Sigma(V). \tag{24}$$

In fact, even when $\Sigma_{\downarrow\downarrow}(V)=\Sigma_{\uparrow\uparrow}(V)=\Sigma(V)$, the MR inversion is provided by another mechanism: the magnon-assisted tunneling. We will now investigate the MR inversion as a consequence of inelastic processes. To model the inelastic transmission coefficients, defined as $T^J_{\sigma,-\sigma}=|\langle R,\sigma|H_I|L,-\sigma\rangle|^2$, we will adopt *Lorentzian* functions of the form

$$T_{\uparrow\downarrow}(\epsilon,\omega) = \rho(\omega)\frac{n(\omega)}{N|\Delta|S}\frac{\Gamma^2_{\uparrow\downarrow}(\epsilon)}{\Gamma^2_{\uparrow\downarrow}(\epsilon) + (\epsilon - \epsilon_{r\uparrow\downarrow}(V))^2},$$

$$T_{\downarrow\uparrow}(\epsilon,\omega) = \rho(\omega)\frac{n(\omega)+1}{N|\Delta|S}\frac{\Gamma^2_{\downarrow\uparrow}(\epsilon)}{\Gamma^2_{\downarrow\uparrow}(\epsilon) + (\epsilon - \epsilon_{r\downarrow\uparrow}(V))^2},$$

where $\rho(\omega)$ is the magnon density of states, $n(\omega) = (\exp[\omega/k_{\mathrm{B}}T]-1)^{-1}$ is the magnon number, following the Bose–Einstein distribution. The appearance of the number of magnons $n(\omega)$ in the transmission coefficients is due to the magnon creation (destruction) operator $b^\dagger_{\mathbf{q}}(b_{\mathbf{q}})$. The magnon energy ω is a consequence of the overall electron–magnon system energy conservation. A magnon can be absorbed (emitted) with consequent spin-flip and an increase (decrease) in the electron's energy. The magnon density of states $\rho(\omega)$ comes from the replacement of an integral in the $\mathbf{q}$ variable to an integral in magnon energy ω. The emission and absorption currents, integrating over all magnon states allowed, are written as follows:

$$I_{\mathrm{em}} = \frac{2\pi e}{\hbar}\int d\epsilon \int d\omega (T_{\downarrow\uparrow'}(\epsilon,\omega)N^L_{\downarrow}(\epsilon - V + \omega)N^R_{\uparrow}(\epsilon) \times f_L(\epsilon - V + \omega)[1 - f_R(\epsilon)] - T_{\downarrow'\uparrow}(\epsilon,\omega) \times N^L_{\uparrow}(\epsilon - V + \omega)N^R_{\downarrow}(\epsilon)f_R(\epsilon + V - \omega)[1 - f_L(\epsilon)]) \tag{25}$$

$$I_{\mathrm{abs}} = \frac{2\pi e}{\hbar}\int d\epsilon \int d\omega (T_{\uparrow\downarrow'}(\epsilon,\omega)N^L_{\uparrow}(\epsilon - V + \omega)N^R_{\downarrow}(\epsilon) \times f_L(\epsilon - V + \omega)[1 - f_R(\epsilon)] - T_{\uparrow'\downarrow}(\epsilon,\omega) \times N^L_{\downarrow}(\epsilon - V + \omega)N^R_{\uparrow}(\epsilon)f_R(\epsilon + V - \omega)[1 - f_L(\epsilon)]). \tag{26}$$

Introducing a low cutoff in the magnon spectrum and taking the low temperature limit $T\to 0$ K, we get $n(\omega)\to 0$ for the Bose–Einstein distribution. This limit excludes the absorption current, leaving only the emission contributions to the total current:

$$I_{\mathrm{em}}^{(\mathrm{P})} = \frac{2\pi e}{N|\Delta|S\hbar}(W_3(V) + W_4(V))\Sigma^1(V), \tag{27}$$

$$I_{\mathrm{em}}^{(\mathrm{AP})} = \frac{2\pi e}{N|\Delta|S\hbar}(W_1(V) + W_2(V))\Sigma^1(V), \tag{28}$$

where we have defined

$$\Sigma^1(V) = \int d\omega \int_0^{V-\omega} d\epsilon\, \rho(\omega)\frac{\Gamma^2_{\downarrow\uparrow}(V)}{\Gamma^2_{\downarrow\uparrow}(V) + (\epsilon - \epsilon_{r\downarrow\uparrow}(V))^2} = \Gamma_{\downarrow\uparrow}(V)\int d\omega\, \rho(\omega)\left[\arctan\left(\frac{V-\omega-\epsilon_{r\downarrow\uparrow}(V)}{\Gamma_{\downarrow\uparrow}(V)}\right) + \arctan\left(\frac{\epsilon_{r\downarrow\uparrow}(V)}{\Gamma_{\downarrow\uparrow}(V)}\right)\right]. \tag{29}$$

When the limit $\Gamma(V)\to 0$ (narrow resonance) is valid, the above equation can be simplified, being the result

$$\Sigma^1(V) = \begin{cases} \pi\Gamma_{\downarrow\uparrow}(V)\int_0^{3V/2-\epsilon_{0m}} d\omega\, \rho(\omega), & \text{for } \frac{2\epsilon_{0m}}{3} < V < 2\epsilon_{0m}; \\ 0, & \text{otherwise}. \end{cases} \tag{30}$$

We put the magnon density of states in a general form:

$$\rho(\omega) = c_1\omega^n \exp[-c_2(\omega - \omega_0)^m].$$

The above expression (when the exponents n and m are adequately chosen) discards the need to introduce a low and a maximum abrupt cutoff in the magnon spectrum. The differential conductance due to magnon-assisted tunneling current is readily obtained through $G_{\mathrm{em}}=dI_{\mathrm{em}}/dV$, being the total differential conductance given by $G=G_d+G_{\mathrm{em}}$. The total currents for both P and AP configurations, including elastic and inelastic terms, are shown below:

$$I^{(\mathrm{P})} = \frac{2\pi e}{\hbar}\left\{[W_1(V) + W_2(V)]\Sigma(V) + \frac{1}{N|\Delta|S}[W_3(V) + W_4(V)]\Sigma^1(V)\right\}, \tag{31}$$

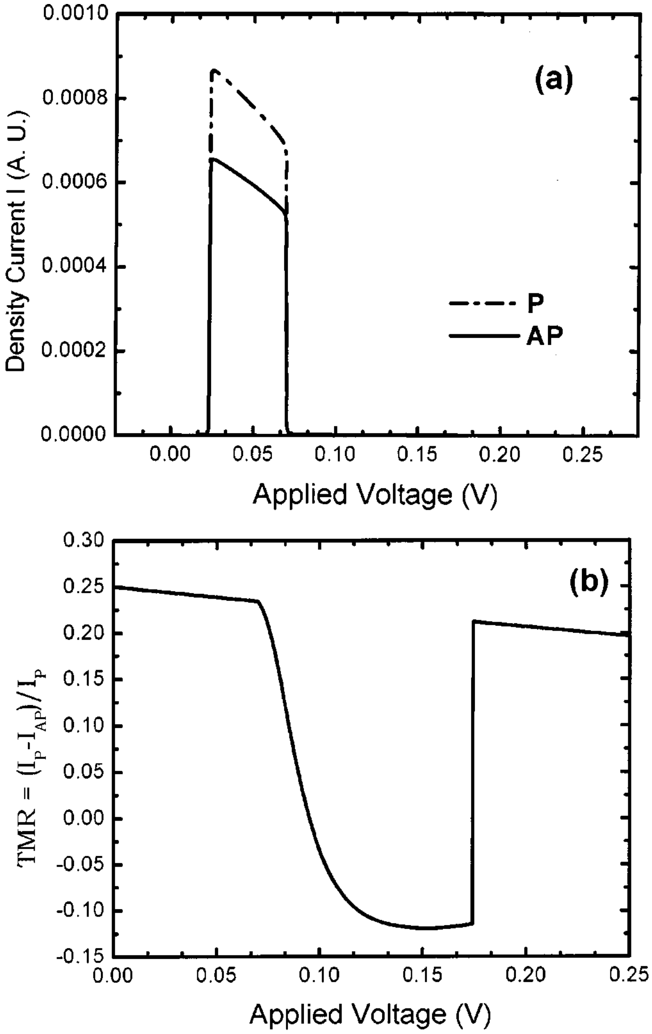

FIG. 3. The MR inversion as a consequence of magnon-assisted tunneling: (a) I-V characteristics and (b) the tunneling magnetoresistance defined by (9), using $\epsilon_{0d}=35$ meV and $\epsilon_{0m}=87.2$ meV.

$$I^{(\mathrm{AP})}=\frac{2\pi e}{\hbar}\left\{[W_3(V)+W_4(V)]\Sigma(V)+\frac{1}{N|\Delta|S}[W_1(V)+W_2(V)]\Sigma^1(V)\right\}. \quad (32)$$

The MR inversion provided by an inelastic mechanism is obtained when the inequality $I^{(\mathrm{P})}(V)<I^{(\mathrm{AP})}(V)$ is satisfied, or, equivalently,

$$\Sigma(V)<\frac{1}{N|\Delta|S}\Sigma^1(V). \quad (33)$$

Since the resonances for direct and magnon-assisted tunneling occur at different electron energies, allowing the condition (33) to be satisfied, the MR inversion is found in a range of voltages $V_1<V<V_2$.

In Fig. 3 the main transport properties of a RMTJ including the magnon-assisted tunneling current are shown. The following set of parameters have been used: $a=0.5$ nm, $V_0=3.0$ eV, $N_m^F=1.0$ in normalized units, $r=2.21$, $\lambda=0.08$, $\beta=2.9$, and $\eta=0.1$. We justify the parameter $N|\Delta|S=3/2$ due to the fact that the product of the exchange energy $|\Delta|$ ($\sim$10–20 meV) with the number of interface atoms N($\sim$20–100 atoms) is of the order of unity. The spin value is also of the order of unity, so we made $S=3/2$. For the magnon density of states, we chose $\rho(\omega)=\omega\exp[-c_2(\omega-\omega_0)^2]$, being $\omega_0=16$ meV and $c_2=500$ eV^{-2}. The value of ω_0 (which corresponds to the peak in the magnon spectrum) is in agreement with the experimental data shown in Ref. 1. Notice that our magnon density of states is not of the form $\sqrt{\omega}$, as one would expect for surface magnons, however eliminating the divergence due to $n(\omega)$ at $\omega=0$, without the need of a lower cutoff in the magnon spectrum. In Fig. 3(a) we show the total currents and in Fig. 3(b) the TMR, calculated with (9). The direct tunneling and inelastic tunneling resonance energies are $\epsilon_{0d}=35$ meV and $\epsilon_{0m}=87.2$ meV, respectively. The RMTJ behaves as expected: the direct resonance occurs first, consequently $I_{\mathrm{P}}>I_{\mathrm{AP}}$ due to the fact that $\Sigma(V)\gg\Sigma^1(V)$ and the band structure favor the parallel configuration $[W_1+W_2>W_3+W_4$, $I_{\mathrm{P}}\propto(W_1+W_2)\Sigma$ and $I_{\mathrm{AP}}\propto(W_3+W_4)\Sigma]$. With the increase of the applied voltage the inelastic process resonance is reached and the direct resonance has been turned off, being this situation being quite different: with $\Sigma(V)\ll\Sigma^1(V)$ the band structure favors the anti-parallel configuration $[W_1+W_2>W_3+W_4$ due to the fact that now $I_{\mathrm{P}}\propto(W_3+W_4)\Sigma^1$ and $I_{\mathrm{AP}}\propto(W_1+W_2)\Sigma^1]$ and, consequently, $I_{\mathrm{P}}<I_{\mathrm{AP}}$. The MR inversion is obtained. For greater voltages, both resonances are turned off; the MR turns out to be positive again and the junction behaves as a simple MTJ. We can change the scenario explained above, introducing other inelastic mechanisms like phonon excitations at larger temperatures and impurity doping of the electrodes. Resonance levels can be provided by impurity scattering centers located inside the insulating barrier of a simple MTJ. However, the assumption that the tunneling amplitude from left to right side tunneling is the same as in the opposite direction is not valid, the transmission amplitudes being dependent on the position of the impurity center inside the barrier.

IV. CONCLUSIONS

We have presented a consistent study of the voltage dependence of the "giant" magnetoresistance in ferromagnetic resonant tunneling junctions. Our approach includes (a) a resonant tunneling coefficient with *Lorentzian* shape and lowering of the effective barrier height with the applied voltage, changing the resonance level; (b) different variations of the density of states for each spin band with voltage; and (c) magnon assisted inelastic tunneling near zero bias, allowing the MR inversion.

The MR inversion is an interesting phenomena that can be provided by different transmission coefficients for spin-up and spin-down channels, with spin-up and spin-down resonances occurring at different energies or through inelastic processes entering via magnon excitation. In the last case the resonance level of the inelastic process is different from the direct tunneling resonance allowing the MR inversion. Temperature effects are not discussed in this paper. As shown in Sec. IV, only magnon emission processes are included at low temperature ($T\to0$). At finite temperature, we expect a decrease of the resistance near zero bias, due to one-magnon-absorption assisted tunneling. The above should be superimposed to the thermal smearing in the Fermi–Dirac distribution of tunneling electrons.[3]

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[1]Y. Ando, J. Murai, H. Kubota, and T. Miyazaki, J. Appl. Phys. **87**, 5209 (2000).
[2]X. H. Xiang, T. Zhu, J. Du, G. Landry, and J. Q. Xiao, Phys. Rev. B **66**, 174407 (2002).
[3]J. J. Akerman, I. V. Roushchin, J. M. Slaughter, R. W. Dave, and I. K. Schuller, Europhys. Lett. **63**, 104 (2003).
[4]F. Montaigne, J. Nassar, A. Vaurs, F. Nguyen Van Dau, F. Petroff, A. Schuhl, and A. Fert, Appl. Phys. Lett. **73**, 2829 (1998).
[5]T. Miyazaki and N. Tezuka, J. Magn. Magn. Mater. **139**, L231 (1995).
[6]G. G. Cabrera and L. M. Falicov, Phys. Status Solidi B **61**, 539 (1974); Phys. Rev. B **11**, 2651 (1975).
[7]M. Jullire, Phys. Lett. **54A**, 225 (1975).
[8]S. Zhang, P. M. Levy, A. C. Marley, and S. S. P. Parkin, Phys. Rev. Lett. **79**, 3744 (1997).
[9]J. S. Moodera, J. Nowak, and R. J. M. van de Veerdonk, Phys. Rev. Lett. **80**, 2941 (1998).
[10]J. S. Moodera and G. Mathon, J. Magn. Magn. Mater. **200**, 248 (1999).
[11]G. G. Cabrera and N. Garcia, Appl. Phys. Lett. **80**, 1782 (2002).
[12]C. A. Dartora and G. G. Cabrera, J. Appl. Phys. **95**, 6058 (2004).
[13]X. Zhang, B. Z. Li, G. Sun, and F. C. Pu, Phys. Rev. B **56**, 5484 (1997); N. Ryzhanova, G. Reiss, F. Kanjouri, and A. Vedyayev, arxiv.cond-mat/0401006 v2, 12 January 2004.
[14]S. Zhang and P. M. Levy, Eur. Phys. J. B **10**, 599 (1999).
[15]E. Y. Tsymbal, A. Sokolov, I. F. Sabirianov, and B. Doudin, Phys. Rev. Lett. **90**, 186602 (2003).
[16]E. Y. Tsymbal and D. G. Pettifor, Phys. Rev. B **64**, 212401 (2001).
[17]R. Jansen and J. S. Moodera, Appl. Phys. Lett. **75**, 400 (1999).
[18]P. LeClair, J. T. Kohlhepp, H. J. M. Swagten, and W. J. M. de Jonge, Phys. Rev. Lett. **86**, 1066 (2001).
[19]J. M. De Teresa *et al.*, Phys. Rev. Lett. **82**, 4288 (1999).
[20]X. Ferry and X. Goodnick, *Transport in Nanostructures* (Cambridge University Press, Cambridge, 1997).
[21]Y. Imry, *Introduction to Mesoscopic Physics* (Oxford University Press, New York, 1997).
[22]C. Kittel, *Quantum Theory of Solids* (Wiley, New York, 1963).