Magnetoresistance through grain boundaries in a resonant-tunneling mechanism

H. Sun,[1,2] K. W. Yu,[1] and Z. Y. Li[2,3]
[1]*Department of Physics, the Chinese University of Hong Kong, Shatin, NT, Hong Kong*
[2]*Department of Physics, Suzhou University, Suzhou 215006, China*
[3]*CCAST (World Laboratory), P.O. Box 8730, Beijing 100080, China*
(Received 16 March 2003; published 19 August 2003)

The effects of local spins inside grain boundaries on tunneling magnetoresistance are studied. A spin-dependent multichannel picture is used to describe the magnetotransport processes across antiferromagnetic grain boundaries. Using a resonant-tunnelling model within a mean-field scheme, we calculate the magnetoresistance caused by the change of the magnetic configuration of grains and inside grain boundaries. We show that, unlike paramagnetic barriers, the existence of local spins in an antiferromagnetic grain-boundary always decreases the low-field magnetoresistance. In the limit of complete spin polarization, a high-field slope proportional to the grain boundary magnetization is found, which is consistent with the characteristics of the high-field magnetoresistance phenomenon. By taking the effects of the direct-, resonant-, and inelastic-tunneling processes into account simultaneously, we can explain the experimental results of extrinsic magnetoresistance in highly spin-polarized magnetic oxides qualitatively.

DOI: 10.1103/PhysRevB.68.054413 PACS number(s): 75.45.+j, 72.25.−b, 73.40.Gk, 45.70.−n

I. INTRODUCTION

Extrinsic magnetoresistance (MR) in grain-boundary systems of metallic oxides with high spin polarization, such as polycrystalline manganite bulk[1] and thin film,[2] CrO_2 pressed powders,[3] etc., has attracted much attention for its scientific and technologic significance.[4] Many models[1,2,5–8] focusing on the relationship between the behavior of the MR and the properties of grain boundaries (GB's) in these systems have been proposed, based on different transport mechanisms but without ruling out each other. Most of them[1,2,5,6] treat the grain boundary as an insulating barrier across which a spin-polarized tunneling occurs between adjacent grains; some[7,8] suggest that it might be a mesoscopic region with distorted magnetic and transport properties. So far none can explain all the observed experimental facts.

Recent nonlinear conductance measurements performed on a lot of the grain-boundary systems of ferromagnetic oxides[6,9–11] gave clear experimental evidence for one of the transport mechanisms: an inelastic-tunneling process between ferromagnetic grains via chains of local states in the GB's. Such a process is characterized by the special power law of the voltage and the temperature dependence of the conductance[12]

$$G=g_s^V V^s, \quad eV \gg k_B T \tag{1}$$

$$G=g_s^T T^s, \quad eV \ll k_B T. \tag{2}$$

Here, $s=N-2/(N+1)$ $(N\geqslant 2)$, and N is the number of the localized states composing the chain. Equations (1) and (2) are the well-known results of the Glazman-Matveev model[12] and were confirmed by Xu *et al.*[13] in the experiments on the conductance of the metal/amorphous silicon/metal junctions. A good agreement between the model and the data from the experiments on a lot of grain-boundary junctions[6,9,10] and some granular systems[11] proved the existence of the inelastic tunneling via local states additional to the direct-tunneling channel in the GB transport. The total magnetoresistance is suppressed due to the spin-polarization loss during the inelastic-tunneling process[4,9,11] and decays rapidly with temperatures or voltages increasing because of the enhanced contribution of the inelastic tunnelling.[9,11]

As Glazman and Matveev pointed out,[12] the presence of localized states in the barrier opens up the possibility of not only the inelastic tunneling described by Eqs. (1) and (2) but also an elastic tunneling via a single impurity in the barrier, which is named as "resonant tunneling" due to its Lorentzian resonant characteristics. At low temperatures the resonant-tunneling process plays a dominant role until a certain temperature is reached. After that, the contribution of inelastic two-impurity channels exceeds the contribution of the resonant tunneling and the inelastic processes of conduction in chains consisting of more impurities are activated with further increasing of temperature. This crossover from elastic resonant tunneling via a single impurity to inelastic tunneling via a chain of local states was also exhibited by experiments.[13] In magnetic junctions having a single barrier with defects[14,15] or in a double-barrier system[16] in which a quasibound state provides the resonant level, resonant tunneling exhibits a particular contribution to the magnetoresistance compared to the direct tunneling . The total magnetoresistance may be reduced or enhanced, depending on detailed environments, and sometimes the effects are expected to be dramatic.[16]

It seems curious that resonant-tunneling processes have not attracted as much attention in the study of GB magnetoresistance. The low-temperature process is usually treated as an "elastic channel" without distinguishing different contributions from the direct tunnelling and the resonant tunneling.[9,11] On the other hand, the effects of local spins inside the barrier on the GB magnetoresistance has been a subject of interest and considered to be responsible for the field dependence of the magetoresistance.[5,17,18]

These considerations motivated the present work on the effects of local spins inside the grain boundaries in a resonant-tunneling mechanism. We adopted a multichannel picture taking into account the interaction between local spins and tunneling electrons, and dealt with the magnetic coupling between local spins using a mean-field approach. The consequent MR ratios, both at low fields and at high fields, are calculated and a high-field slope of magnetoconductance is obtained. The results show a typical field dependence of the magnetoresistance found in the considered materials. We expect that, after considering the GB properties in a more detailed way, most of the experimental results of the magnetotransport properties of the grain-boundary systems will be explained by a model combining the resonant-tunneling process with the direct tunneling and higher-order inelastic tunneling.

II. RESONANT-TUNNELING MODELS

As mentioned above, the resonant tunneling in a single-barrier tunnel junction refers to the electron tunneling via one impurity inside the barrier. The conductance due to an elastic electron transition from the left electrode to the right via a localized state with energy ϵ_i at the site $\mathbf{c}_i$ in a barrier of the height ϕ and width w is given by[19]

$$g(\epsilon_i,\mathbf{c}_i)=\frac{e^2}{2\pi\hbar}\frac{4\Gamma_L\Gamma_R}{(\epsilon-\epsilon_i)^2+\Gamma^2} \tag{3}$$

where $\Gamma=\Gamma_L+\Gamma_R$ is the inverse lifetime of the resonant states, given by a sum of the partial widths $\Gamma_{L(R)}$ corresponding to electron tunneling from the local state to the left (right) electrode,[20]

$$\Gamma_{L(R)}=\frac{2mq_{L(R)}}{m^2k_0^2+q_{L(R)}^2}\frac{e^{-2k_0w_{L(R)}}}{k_0w_{L(R)}}. \tag{4}$$

$w_{L(R)}$ is the distance of the local state from the left (right) electrode, $q_{L(R)}$ is the Fermi wave vector in the left (right) electrode, m is the effective electron mass in the barrier, and $k_0=\sqrt{2m\phi/\hbar^2}$. Assuming a uniform distribution of the localized states in energy space with a density ν near the Fermi level, the resonant-tunneling conductance via an impurity at $\mathbf{c}_i$ integrated over the energy space can be given in a good approximation by[15,21]

$$g(\mathbf{c}_i)=\frac{2e^2\nu}{\hbar}\frac{\Gamma_L\Gamma_R}{\Gamma_L+\Gamma_R}. \tag{5}$$

If the electrodes are ferromagnetic, the exchange splitting of the energy bands of the electrodes will lead to the formation of two spin-dependent conductive channels [see Fig. 1(b)], one for spin-up tunneling electrons and the other for spin-down (the spin direction is with respect to the magnetization orientation of the grains), and consequently a magnetoresistance due to the resonant tunneling. The conductance for an electron with spin σ is determined by[20]

$$g_\sigma(\mathbf{c}_i,\epsilon_i)=\frac{e^2}{2\pi\hbar}\frac{4\Gamma_{L\sigma}\Gamma_{R\sigma}}{(\epsilon-\epsilon_{i\sigma})^2+\Gamma^2} \tag{6}$$

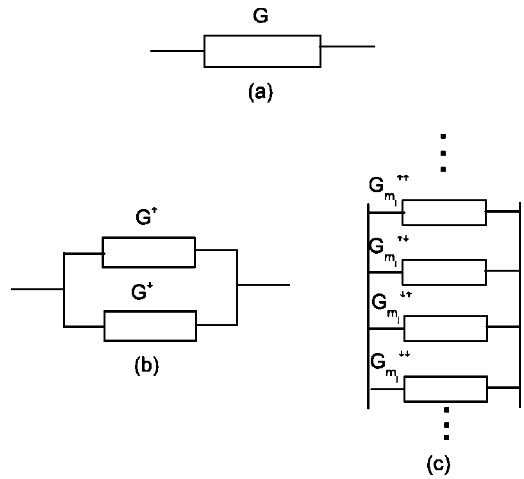

FIG. 1. The channel illustrations for different resonant-tunneling processes: (1) monochannel in nonmagnetic junctions with nonmagnetic impurities; (2) a two-channel process in ferromagnetic junctions with nonmagnetic impurities; (3) a multichannel process in ferromagnetic junctions with paramagnetic impurities

with

$$\Gamma_{L\sigma(R\sigma)}=\frac{2mq_{L\sigma(R\sigma)}}{m^2k_0^2+q_{L\sigma(R\sigma)}^2}\frac{e^{-2k_0w_{L(R)}}}{k_0w_{L(R)}}. \tag{7}$$

In a more complicated case of ferromagnetic electrodes with paramagnetic impurities inside the barrier, the strong exchange interaction between the local spin and the tunneling electron will lead to an increase of the number of channels as well as a modification to the resonant levels and consequently the correspondent inverse lifetime, as shown in the theoretical work of Vedyayev *et al.*,[21] where the spin-dependent resonant tunneling across a barrier was studied with paramagnetic impurities at an effective field from the exchange interaction with the ferromagnetic electrodes. Considering a tunneling process via a local spin $\mathbf{S}$ at an effective field $\mathbf{H}_{eff}$, a multichannel picture of resonant tunneling deduced from their calculation results can be described as follows.

(1) Supposing a local spin with an initial state $|S,m\rangle$, there are four channels according to the initial and final states of the tunneling electron and the local spin [see Fig. 1(c)]. The spin of the tunneling electron invaries in two of the channels and flips in the other two while the total magnetic component of the spins of the electron and the impurity keeps conserved. So the channels can be labeled by μ, ρ, and m_j, where μ, ρ are the electron spins before and after the transition, respectively, and $m_j=\mu+m$ is the total spin component along the field direction.

(2) Suppose the scattering potential amplitude on the impurity is Δ_i and the exchange interaction between the electron and the local spin is J. The possible resonant state is assumed to correspond to the multiplet $\epsilon_i=\Delta_i-(S/2)J$ with a total angular momentum $j=S+\frac{1}{2}$. The contribution of the other multiplet $\epsilon_i'=\Delta_i+[(S+1)/2]J$ can be neglected since ϵ_i is much closer to the Fermi energy than the latter because

of the large value (~ 1 eV) of J and the strong Coulomb repulsion on impurities in the considered materials.[18,21] It means only the state

$$|j=S+\tfrac{1}{2},m_j\rangle = C_{1/2}|S,m=m_j-\tfrac{1}{2};\tfrac{1}{2},\mu=\uparrow\rangle + C_{-1/2}|S,m=m_j+\tfrac{1}{2};\tfrac{1}{2},\mu=\downarrow\rangle \tag{8}$$

goes into resonance. Then the correspondent inverse lifetime, which is related to the possibilities of electron tunneling from the local state to the left and right electrodes,[20] can be given by

$$\gamma = (C_{1/2}^2\Gamma_{L\uparrow} + C_{-1/2}^2\Gamma_{L\downarrow}) + (C_{1/2}^2\Gamma_{R\uparrow} + C_{-1/2}^2\Gamma_{R\downarrow}) = C_{1/2}^2\Gamma_{\uparrow} + C_{-1/2}^2\Gamma_{\downarrow} \tag{9}$$

with $C_{\pm 1/2}$ denoting the Clebsch-Gordan coefficients,

$$C_{\pm 1/2} = \sqrt{\frac{S \pm m_j + \frac{1}{2}}{2S+1}}, \tag{10}$$

and the partial widths for the left and right electrodes are, respectively, given by

$$\gamma_L^{\mu} = C_{\mu}^2 \Gamma_{L\mu}, \tag{11}$$

$$\gamma_R^{\rho} = C_{\rho}^2 \Gamma_{R\rho}. \tag{12}$$

Here $\Gamma_{L\sigma}(\sigma=\uparrow,\downarrow)$ is determined by Eq. (7) and $\Gamma_{\sigma}=\Gamma_{L\sigma}+\Gamma_{R\sigma}$.

(3) The statistical probabilities $P_{m_j}^{\mu\rho}$ of a channel (m_j,μ,ρ), obtained from the calculated results in Ref. 21, are adjusted by the temperature-reduced effective field $h=\mu_B H_{eff}/k_B T$. For spin-conserved channels ($\mu=\rho$),

$$P_{m_j}^{\mu\mu} = P(S_z = m_j - \mu) = \frac{\exp(2S_z h)}{\sum_{m=-S}^{S}\exp(2mh)}, \tag{13}$$

which is also the statistical probability of the z component of local spins at the effective field, while for spin-flipping channels,

$$P_{m_j}^{\uparrow\downarrow} = P_{m_j}^{\downarrow\uparrow} = \frac{1}{2}P\left(S_z = m_j - \frac{1}{2}\right)\int f_-(\epsilon)d\epsilon + \frac{1}{2}P\left(S_z = m_j + \frac{1}{2}\right)\int f_+(\epsilon)d\epsilon, \tag{14}$$

where

$$f_-(\epsilon) = \frac{1}{k_B T} f_{\uparrow}(1-f_{\downarrow})$$

$$f_+(\epsilon) = \frac{1}{k_B T} f_{\downarrow}(1-f_{\uparrow})$$

and

$$f_{\uparrow,\downarrow} = \frac{1}{1+e^{(\epsilon-\epsilon_F \mp \mu_B H_{eff})/k_B T}}. \tag{15}$$

Then the total conductance for a channel (m_j,μ,ρ) is given by the combination of the quantum-mechanical probability $\sigma_{m_j}^{\mu\rho}$ and the statistical probability $P_{m_j}^{\mu\rho}$

$$g_{m_j}^{\mu\rho} = g_0 \sigma_{m_j}^{\mu\rho} P_{m_j}^{\mu\rho}, \tag{16}$$

where

$$\sigma_{m_j}^{\mu\rho} = \frac{\gamma_L \gamma_R}{\gamma} \tag{17}$$

is obtained by substituting the results of the tunnel widths into the expression of the resonant-tunneling conductance as shown in Eq. (5), and g_0 is the independent coefficient including the quantum of conductance and the density ν of the distribution of the resonant levels.

III. GB TMR

Now we begin to study the spin-dependent tunneling through a grain boundary in highly spin polarized magnetic oxides such as perovskite manganite. Unlike tunnel junctions with an amorphous[13] or δ-doped[14,15] barrier, the grain-boundary region contains lots of local spins interacting with each other.[4] Here we consider the case of a barrier with antiferromagnetically coupled local spins, which usually models the GB region for typical half-metallic oxides manganite[4,18] and CrO_2.[22] To calculate transport processes via these local spins is actually a many-body problem,[18] but the resonant-tunneling characteristics exhibited in experiments[9–11] make us believe it is acceptable to deal with the magnetic coupling in a mean-field approximation so that the total conductance can be expressed by the sum of all resonant-tunneling processes via a local spin $\mathbf{S}$ at an effective field

$$G = \sum_i g(\mathbf{c}_i, \mathbf{H}_m^i) = \sum_{i\in A} g(\mathbf{c}_i, \mathbf{H}_m^A) + \sum_{i\in B} g(\mathbf{c}_i, \mathbf{H}_m^B). \tag{18}$$

Here, $g(\mathbf{c_i},\mathbf{H}_m^{\alpha})$ ($\alpha=A,B$) is the conductance via a local spin at the site $\mathbf{c}_i$ on sublattice α, and $\mathbf{H}_m^{\alpha}$ is the conventional effective field for sublattice α in antiferromagnets,

$$\mathbf{H}_m^A = \mathbf{H}_0 + \lambda \mathbf{M}_B, \tag{19}$$

$$\mathbf{H}_m^B = \mathbf{H}_0 + \lambda \mathbf{M}_A, \tag{20}$$

with $\mathbf{H}_0$ denoting the applied field, $\mathbf{M}_{\alpha} \equiv \langle S_z^{\alpha}\rangle$ is the average value of the z component of the local spins on sublattice α, and λ is the mean-field parameter characterizing the exchange interaction between the sublattices. The intrasublattice interactions have been neglected.

Substituting the expressions of the resonant-tunneling conductance in Sec. II into Eq. (17) and considering that only the impurities located within k_0^{-1} of the center of the barrier contribute most to the conductance,[13] the total con-

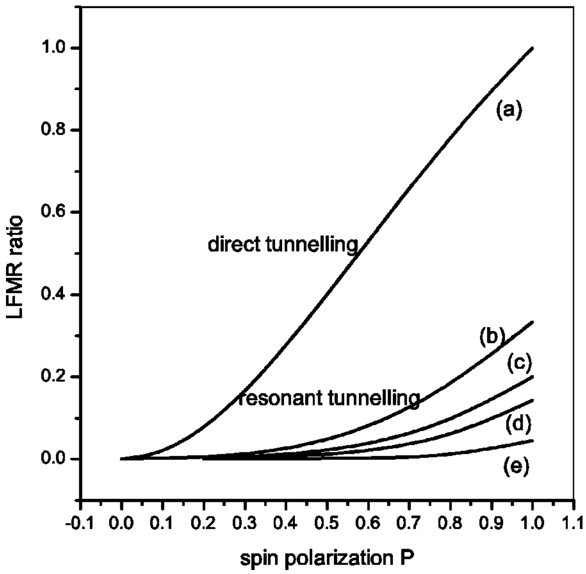

FIG. 2. The LFMR ratio vs the spin polarization P for (a) direct-tunneling processes and resonant-tunneling processes via local spins with (b) $S=1/2$, (c) $S=3/2$, (d) $S=5/2$, and (e) $S=10$.

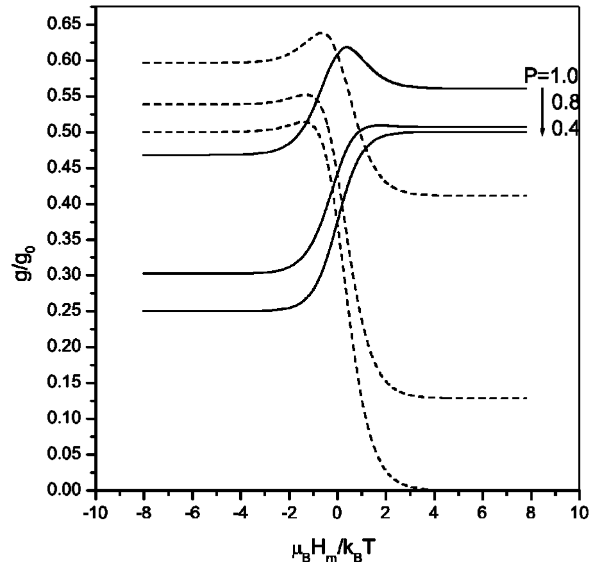

FIG. 3. The calculated results of the normalized resonant-tunneling conductance, g/g_0, via a local spin $S=1/2$ at the center of the barrier for different values of the electrode spin polarization P. The solid lines represent the case of the parallel alignment of the electrode magnetization, and the dashed lines represent the antiparallel alignment.

ductance across a barrier with antiferromagnetically coupled spins inside it can be given by

$$G=G_0\left[\sum_{m_j,\mu,\rho}\sigma_{m_j}^{\mu\rho}\left(c_i=\frac{w}{2}\right)P_{m_j}^{\mu\rho}(H_m^A)+\sum_{m_j,\mu,\rho}\sigma_{m_j}^{\mu\rho}\left(c_i=\frac{w}{2}\right)P_{m_j}^{\mu\rho}(H_m^B)\right]. \tag{21}$$

G_0 is a constant from the summation over the distribution of resonant levels in energy and position space and will be canceled out in the following calculation.

A. Low-field magnetoresistance

First we calculate the magnetoresistance caused by the change of the relative magnetic orientation of the electrodes, which corresponds to the large and rapid decrease of the resistance observed in GB systems of magnetic oxides, i.e., the low-field magnetoresistance (LFMR).[1–4] Compared with the high-field region in which the resistance decreases slowly, the magnetic-field scale of the low-field magnetoresistance is much smaller and has much less effect on the spin configuration inside the barrier. So we can assume only an infinitesimally small field is needed to change the initial antiparallel (AP) alignment of magnetization of the ferromagnetic electrodes to a parallel (P) alignment,[18] and the MR ratio defined as

$$MR=\frac{G_P-G_{AP}}{G_P} \tag{22}$$

can be calculated using Eq. (20) by taking $H_0=0$.

Figure 2 shows the dependence of the MR ratio on the spin polarization of the electrodes

$$P\equiv\frac{q_\uparrow-q_\downarrow}{q_\uparrow+q_\downarrow} \tag{23}$$

for different angular momenta S of local spins at low temperature $T=4.2$ K, and compares it with the results of direct tunneling using the Julliere model[23]

$$MR^{dir}=\frac{2P^2}{1+P^2}. \tag{24}$$

Fixed parameters are set as $q_\uparrow=1$ Å^{-1}, $k_0=0.5$ Å^{-1}, $m=0.4$ according to the typical values in ferromagnetic junctions,[21] and the mean-field parameter $\lambda=100$ K. It can be seen that the MR ratio due to resonant tunneling is smaller than that due to the direct tunneling in the entire range of the spin polarization, and the deviation is enlarged with S increasing. At low temperatures, direct and resonant processes coexist and the value of the total magnetoresistance depends on their relative proportions. For perovskite manganites $S=\frac{3}{2}$, the magnetoresistance due to resonant-tunneling processes is only about 20% even in the limit of complete spin polarization $P=1$, so the total LFMR is expected to vary between 20% and 100%, which is in agreement with the observations in grain-boundary junctions.[4]

Thus the existence of local spins in the present case always affects the total magnetoresistance in a negative way, which is different from the results in the case of paramagnetic impurities where the magnetoresistance may be either enhanced or suppressed.[14,15,17,21] Such a difference mainly results from the existence of the two sublattices of local spins because of the antiferromagnetic coupling. As Fig. 3 shows, although the local spins on sublattice A ($H_{eff}>0$) can lead to a dramatic MR when the spin polarization is

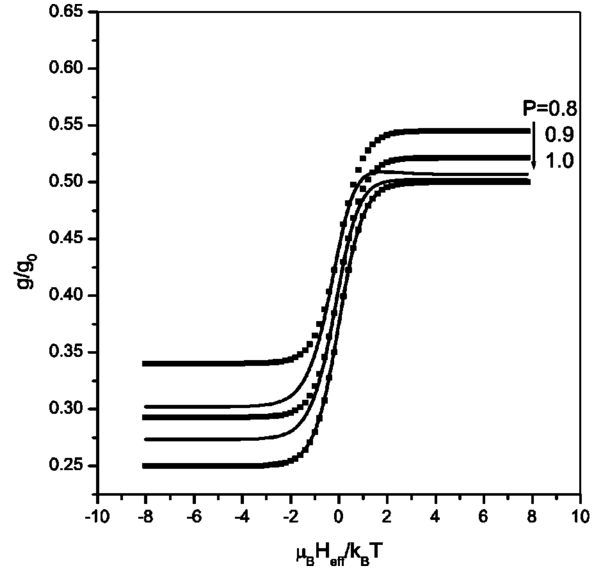

FIG. 4. The dependence of the normalized conductance g/g_0 due to resonant tunneling on the effective field $\mu_B H_{eff}/k_B T$ from exact calculation (lines) and the approximation Eq. (24) (squares) for high spin polarizations.

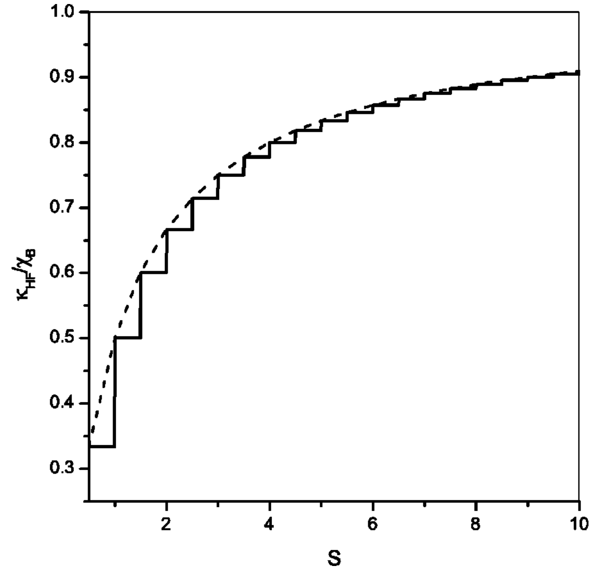

FIG. 5. The dependence of the ratio between the high-field slope of the conductance κ_{HF} and the GB susceptibility χ_B on local spins S in the limit of half-metallicity.

large, the local spins on sublattice B ($H_{eff}<0$) will produce opposite effects and lead to a decrease of the total MR.

B. High-field magnetoresistance

After the P configuration of magnetization of the electrodes is reached at a small field, the distribution of the local spins inside the barrier keeps changing with the applied field increasing further. From Eq. (20) we can see that it will change the probability $P_{m_j}^{\mu\rho}$ of the channels and the total conductance consequently. This magnetoresistance is a "high-field magnetoresistance, which refers to the gradual and linearlike decrease of the resistance following the dramatic and rapid LFMR observed in many GB systems.[4] The high-field magnetoresistance (or magnetoconductance, as defined in Ref. 5) phenomenon is usually described by a normalized high-field slope of the conductance $\kappa_{HF}\equiv 1/[G(H0=0)](dG/dH_0)$.

In order to clarify the complicated dependence of the resonant-tunneling conductance on the applied field, we adopted a simplifying assumption for the high-spin-polarization systems investigated:

$$P_{m_j}^{\mu\rho}=\tfrac{1}{2}[P(S_z=m_j-\tfrac{1}{2})+P(S_z=m_j-\tfrac{1}{2})](\mu\neq\rho). \quad (25)$$

Comparing it with the exact expression shown in Eq. (13), we can see that this assumption neglects the effects of the effective field on the Fermi distribution of the electrons in spin-flipping channels. When the spin polarization is high, the spin-flipping channels contribute little with a small number of the spin-down electrons and the approximation become acceptable. Figure 4 illustrates the curves of the conductance vs the effective field for different spin polarizations calculated using the exact and the approximate expression, respectively. The approximation leads to a deviation from the exact results, but the characteristics of the field-dependence of the resonant-tunneling conductance are reserved. The deviation gradually vanished with the spin polarization approaching the limit of the half-metallicity ($P=1$). Substituting Eq. (24) into the conductance expression (20) and rearranging the results, we get

$$G\propto\frac{S+1}{2(2S+1)}(\Gamma_\uparrow+\Gamma_\downarrow)+\frac{S}{2(2S+1)}(\Gamma_\uparrow-\Gamma_\downarrow)M_B, \quad (26)$$

where $M_B\equiv\langle S_z\rangle/S$ is the average normalized magnetization of local spins. From this linear dependence of the conductance on the average magnetization, we obtain the high-field slope of the conductance due to resonant tunneling via local spins

$$\kappa_{HF}=\frac{S}{S+1}P'\chi_B \quad (27)$$

with

$$P'=\frac{\Gamma_\uparrow-\Gamma_\downarrow}{\Gamma_\uparrow+\Gamma_\downarrow} \quad (28)$$

which represents the unbalance of the two channels without considering the interaction between the local spins and the conducting electrons and is usually increased with the spin polarization increasing. $\chi_B\equiv\partial M_B/\partial H_0$ is the spin susceptibility of the grain boundary. The variation of the proportionality constant κ_{HF}/χ_B with the local spins S for half-metallic systems is shown in Fig. 5.

The linear relationship between the high-field slope and the GB spin susceptibility for high-spin-polarization systems reveals the effects of the spin configuration inside the barrier on the tunneling conductance. Note that this relationship seems to have been a consensus although direct experimental evidences are still unavailable.[4] It is interesting that many models taking into account the effects of GBs, no matter how different mechanisms they are based on, drew the same conclusion that the high-field slope is proportional to the GB susceptibility.[5,8,17] The experiments detecting the GB magnetism according to this proportional relation obtained the results coincident with those from other independent measurements.[4] The agreement of the results of the present model with this well-known prediction can be considered as another supportive proof for the existence of resonant-tunneling processes.

IV. CONCLUSIONS

From all the results we gained a complete picture of how the resonant-tunneling processes through an antiferromagnetic barrier acted on the field behavior of magnetoresistance and concluded that the resonant-tunneling processes make a particular contribution to the magnetoresistance. Thus the grain-boundary conductance should be divided into three parts:

$$G = G^{dir} + G^{res} + G^{inel}, \tag{29}$$

where G^{dir}, G^{res}, and G^{inel} represent the direct, the resonant, and the higher-order inelastic-tunneling conductance, respectively. Note that this model is different from the three-current theory proposed by Höfener *et al.* in which the direct and the resonant processes are considered as one elastic channel.[9] Here, the direct-tunneling conductance G^{dir} provides a background low-field magnetoresistance with the Julliere model as its simplest description. The resonant-tunneling part G^{res} causes both magnetoresistance decreasing at low fields and an extra high-field magnetoresistance proportional to the grain-boundary magnetization. The spin-independent inelastic-tunneling channel G^{inel} via chains of local states is detrimental to the total magnetoresistance in the entire range of the applied field. With temperatures or voltages increasing, both the low-field and the high-field magnetoresistance decrease because of the increasing contribution of G^{inel} and, because of its larger initial values, the low-field magnetoresistance will drop much more rapid by than the latter, which is in agreement with the experimental results.[4,5]

In summary, we calculated the grain-boundary magnetoresistance at low and high fields on the basis of a spin-dependent resonant-tunneling model. The results confirmed further the existence of the resonant-tunneling mechanism which was implied by the observed higher-order inelastic-tunneling characteristics in experiments. By using a general transport picture incorporating direct-, resonant-, and inelastic-tunneling processes, many experimental results can be explained qualitatively, including the following.

(1) The measured values of low-field magnetoresistance at low temperatures are usually smaller than that expected in direct-tunneling models.

(2) The magnetoresistances at low fields and high fields exhibit obviously different behaviors.

(3) The I-V (or G-T) curves have special nonlinear characteristics.

(4) The magnetoresistance will decay rapidly with temperature or voltage increase.

Of course, the conditions of real grain boundaries are much more complex than what we have considered here. For example, they are very likely to show some spin-glass-like properties due to the competition between the ferromagnetic double-exchange and antiferromagnetic superexchange interaction.[4] Also the thickness of the grain-boundary region, which naturally has a distribution in granular systems, will affect the proportions of the three transport mechanisms.[13] We expect, after considering more details of grain boundaries, that a more accurate description based on the present model should be obtained of the transport mechanisms in grain-boundary systems of highly spin-polarized metallic oxides.

ACKNOWLEDGMENTS

This work was supported in part by the Direct Grant for Research and in part by National Natural Science Foundation of China under Grant No. 10174049.

[1] H.Y. Hwang, S.W. Cheong, N.P. Ong, and B. Batlogg, Phys. Rev. Lett. **77**, 2041 (1996).

[2] A. Gupta, G.Q. Gong, G. Xiao, P.R. Duncombe, P. Lecoeur, P. Trouilloud, Y.Y. Wang, V.P. Dravid, and J.Z. Sun, Phys. Rev. B **54**, R15 629 (1996).

[3] J.M.D. Coey, A.E. Berkowitz, Ll. Balcells, F.F. Putris, and A. Barry, Phys. Rev. Lett. **80**, 3815 (1998).

[4] M. Ziese, Rep. Prog. Phys. **65**, 143 (2002), and the references therein.

[5] S. Lee, H.Y. Hwang, B.I. Shraiman, W.D. II Ratcliff, and S.-W. Cheong, Phys. Rev. Lett. **82**, 4508 (1999).

[6] M. Ziese, Phys. Rev. B **60**, R738 (1999).

[7] N.D. Mathur, G. Burnell, S.P. Isaac, T.J. Jackson, B.-S. Teo, J.L. MacManus-Driscoll, and L.F. Cohen, Nature (London) **387**, 266 (1997).

[8] J.E. Evetts, M.G. Blamire, N.D. Mathur, S.P. Isaac, B.-S. Teo, L.F. Cohen, and J.L. MacManus-Driscoll, Philos. Trans. R. Soc. London, Ser. A **356**, 1593 (1998).

[9] C. Höfener, J.B. Philipp, J. Klein, L. Alff, A. Marx, B. Büchner, and R. Gross, Europhys. Lett. **50**, 681 (2000).

[10] N. Khare, U.P. Moharil, A.K. Gupta, A.K. Raychaudhuri, S.P. Pai, and R. Pinto, Appl. Phys. Lett. **81**, 325 (2002).

[11] J. Dai and J. Tang, Phys. Rev. B **63**, 064410 (2001).

[12] L.I. Glazman and K.A. Matveev, Zh. Éksp. Teor. Fiz. **94**, 232 (1988) [Sov. Phys. JETP **67**, 1276 (1988)].

[13] Y. Xu, D. Ephron, and M.R. Beasley, Phys. Rev. B **52**, 2843 (1995).

[14] R. Jansen and J.S. Moodera, Appl. Phys. Lett. **75**, 400 (1999).

[15] R. Jansen and J.C. Lodder, Phys. Rev. B **61**, 5860 (2000).

[16] A.G. Petukhov, A.N. Chantis, and D.O. Demchenko, Phys. Rev.

Lett. **89**, 107205 (2002).
[17] F. Guinea, Phys. Rev. B **58**, 9212 (1998).
[18] J. Inoue, N. Nishimura, and H. Itoh, Phys. Rev. B **65**, 104433 (2002).
[19] A.I. Larkin and K.A. Matveev, Zh. Éksp. Teor. Fiz. **93**, 1030 (1987) [Sov. Phys. JETP **66**, 580 (1987)].
[20] A.M. Bratkovsky, Phys. Rev. B **56**, 2344 (1997).
[21] A. Vedyayev, D. Bagrets, A. Bagrets, and B. Dieny, Phys. Rev. B **63**, 064429 (2001).
[22] Jianbiao Dai and Jinke Tang, Appl. Phys. Lett. **77**, 2840 (2000).
[23] M. Julliere, Phys. Lett. **54A**, 225 (1975).