Theory of the temperature-dependent giant magnetoresistance in magnetic multilayers

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The giant magnetoresistance (GMR) in magnetic multilayers at finite temperatures is studied with the use of the functional-integral method in the static approximation. Numerical calculations are performed by using a simple model. The effect of spin fluctuations is shown to play important roles in discussing the temperature-dependent GMR. The temperature dependence of the calculated GMR is much larger than that of the magnetization, and it shows an almost linear decrease near the Curie temperature when the temperature is raised. It is also shown that the temperature dependence of the GMR is more significant in a multilayer with a larger ground-state GMR. Our calculation accounts well for the GMR feature observed in many transition-metal multilayers.

I. INTRODUCTION

During the last few years, the giant magnetoresistance[1] (GMR) and related phenomena[2] have been intensively investigated in many magnetic multilayers. Theoretically, the ground-state GMR has been studied with a semiclassical approach[3,4] using the Boltzman equation or a microscopic approach[5–8] based on the Kubo formula. The GMR at *finite temperatures* has been discussed so far using the local-spin model.[9] Since most magnetic multilayers consist of transition metals, it is desirable to study their GMR based on the itinerant-electron model.

It was a long-standing problem whether various physical properties of transition metals can be described by the local-spin (Heisenberg) model or the itinerant-electron (band) model. The Curie-Weiss susceptibility and the large specific heat near the Curie temperature are explained by the former model while the nonintegral ground-state moment, the large specific-heat coefficient, and the high conductivity favor the latter model. It has been realized[10] that these two aspects can be reconciled if we take into account the effect of spin fluctuations, which is neglected in the conventional Hartree-Fock approximation to the itinerant model. Hasegawa[11] and Hubbard[12] independently proposed a finite-temperature band theory for bulk magnetism with the use of the functional-integral method within the static approximation. It is a mean-field theory in which the system is regarded as a collection of local magnetic moments. This approach has proved useful in understanding the magnetic, thermodynamical, and transport properties of transition metals at finite temperatures. In particular, it has been shown[13] from a comparison with Monte Carlo simulations that our theory well reproduces the U-T phase diagram of the simple-cubic[14] and infinite-dimensional[15] Hubbard models (U: interaction, T: temperature). This method has been successfully extended and applied to transition-metal alloys[16] and multilayers.[17]

Quite recently, the present author[8] derived a simple, analytic expression for the conductivity of magnetic films using the coherent potential approximation (CPA).[18] The conductivity for currents parallel to the interface is given in terms of the coherent potential (one-electron self-energy) of film layers. By employing the expression obtained, we have discussed the ground-state GMR of magnetic multilayer.[8]

The purpose of the present paper is to incorporate the effect of spin fluctuations into the calculation of the conductivity of magnetic multilayers in order to discuss their GMR at finite temperatures. The paper is organized as follows: In Sec. II we discuss the calculation method of the conductivity and GMR, after briefly reviewing our spin-fluctuation theory.[11] Numerical calculations using a simple model are presented in Sec. III. Section IV is devoted to conclusion and supplementary discussion.

II. CALCULATION METHOD

A. The spin fluctuation theory in the static approximation

We assume an N_f-layer thin film consisting of magnetic A and nonmagnetic B atoms with the simple-cubic (001) interface. The layer parallel to the interface is assigned by the index $n(=1-N_f)$. We assume that A and B atoms are randomly distributed on layer n with the concentrations of x_n and y_n, respectively $(x_n+y_n=1)$. For a given thin film, we adopt the single-band Hubbard model given by

$$H=\sum_s\sum_j \varepsilon_j c_{js}^{\dagger}c_{js}+\sum_s\sum_{jl}t_{jl}c_{js}^{\dagger}c_{ls}+\sum_j U_j n_{j\uparrow}n_{j\downarrow}, \tag{1}$$

where c_{js} is an annihilation operator of an electron with spin $s(=\uparrow,\downarrow)$ on the lattice site j, $n_{js}=c_{js}^{\dagger}c_{js}$, and t_{jl} is the hopping integral. The atomic potential, ε_j, and the on-site interaction, U_j, are assumed to be given by ε^{λ} and U^{λ} when the lattice site j is occupied by a λ $(=A,B)$ atom.

For a study of the finite-temperature band magnetism of the magnetic film, we adopt the single-site functional-integral method proposed by the present author.[11] When we apply the functional-integral method to the model Hamiltonian given by Eq. (1), the partition function is given within the static approximation by[11,12]

$$Z = \int \prod_j d\nu_j \prod_j d\zeta_j \, \exp[-\beta(\phi_0 + \phi_1)], \tag{2}$$

with

$$\phi_0 = (1/4) \sum_j U_j(\nu_j^2 + \zeta_j^2), \tag{3}$$

$$\exp(-\beta\phi_1) = \mathrm{Tr}\, \exp(-\beta H_{\mathrm{eff}}), \tag{4}$$

$$H_{\mathrm{eff}} = \sum_s \sum_j [(\varepsilon_j - (i/2)U_j\nu_j)\, n_j - (1/2)sU_j\zeta_j\, m_j] + H_0^t. \tag{5}$$

Here $n_j = n_{j\uparrow} + n_{j\downarrow}$, $m_j = n_{j\uparrow} - n_{j\downarrow}$, and H_0^t denotes the second (hopping) term in Eq. (1). Equations (2)–(5) show that the partition function can be evaluated by calculating the partition function of the effective one-electron system given by H_{eff} including the random charge (ν_j) and exchange (ζ_j) fields. The former field is included by the saddle-point approximation and the latter field by the alloy-analogy approximation with the CPA. By using the decoupling approximation, we modified the CPA equation as given by[16]

$$x_n\, T_{ns}^A + y_n\, T_{ns}^B = 0, \tag{6}$$

with

$$T_{ns}^\lambda = \frac{\tilde{\varepsilon}_n^\lambda - s(U_n^\lambda/2)\langle\zeta_n^\lambda\rangle - \Sigma_{ns} + [(U_n^\lambda/2)^2\langle(\zeta_n^\lambda)^2\rangle - (\Sigma_{ns} - \tilde{\varepsilon}_n^\lambda)^2]F_{ns}}{[1 - (\tilde{\varepsilon}_n^\lambda - \Sigma_{ns})F_{ns}]^2 - (U_n^\lambda/2)^2\langle(\zeta_n^\lambda)^2\rangle}, \tag{7}$$

where $\tilde{\varepsilon}_n^\lambda = \varepsilon_n^\lambda + (U_n^\lambda/2)\langle N_n^\lambda\rangle$, the coherent potential for an s-spin electron on the layer n, Σ_{ns}, is a function of ε_n^λ, $\langle\zeta_n^\lambda\rangle$, $\langle(\zeta_n^\lambda)^2\rangle$ and $\langle N_n^\lambda\rangle$ $(= -i\langle\nu_n^\lambda\rangle)$. The self-consistent equations for $\langle\zeta_n^\lambda\rangle$, $\langle(\zeta_n^\lambda)^2\rangle$, and $\langle N_n^\lambda\rangle$ are given by

$$\langle\zeta_n^\lambda\rangle = \int d\zeta_j\, \zeta_j\, C_n^\lambda(\zeta_j), \tag{8}$$

$$\langle(\zeta_n^\lambda)^2\rangle = \int d\zeta_j\, \zeta_j^2\, C_n^\lambda(\zeta_j), \tag{9}$$

$$\langle N_n^\lambda\rangle = \int d\varepsilon\, f(\varepsilon) \sum_s (-1/\pi)\, \mathrm{Im} F_{ns}^\lambda(\varepsilon), \tag{10}$$

where $f(\varepsilon)$ is the Fermi-distribution function. We should note that $F_{ns}^\lambda(\varepsilon)$, the local Green function of an s-spin electron at a λ atom on layer n, and $C_n^\lambda(\zeta_j)$, the distribution of the potential of $(U_n^\lambda/2)\zeta_j$ when a λ atom occupies the layer n, are functions of the coherent potentials, Σ_{ns}, which are functions of $\langle\zeta_n^\lambda\rangle$, $\langle(\zeta_n^\lambda)^2\rangle$, and $\langle N_n^\lambda\rangle$ [Eqs. (6) and (7)]. Thus these quantities must be solved simultaneously, details having been reported elsewhere.[11,16,17]

Once these are determined, we can obtain various physical quantities. For example, the average of the magnetic moment and its root-mean-square (rms) values of a λ atom on the n layer are given by

$$\langle M_n^\lambda\rangle = \langle\zeta_n^\lambda\rangle, \tag{11}$$

$$\langle(M_n^\lambda)^2\rangle^{1/2} = [\langle(\zeta_n^\lambda)^2\rangle - (2T/U_n^\lambda)]^{1/2}. \tag{12}$$

B. Conductivity and GMR

It has been shown that the conductivity for currents parallel to the film layer is given, within the Born approximation, by[8]

$$\sigma = N_f^{-1} \sum_n \sigma_n, \tag{13}$$

with

$$\sigma_n = (e/\hbar)^2\pi \sum_s \sum_l \int d\varepsilon \left(-\frac{\partial f}{\partial\varepsilon}\right) \frac{\nu_s a_{nls}\, \tau_{nls}}{(\Delta_{ns} + \Delta_{ls})}, \tag{14}$$

$$\tau_{nls} = \delta_{nl} + (1 - \delta_{nl}) \left(\frac{(\Delta_{ns} + \Delta_{ls})^2}{[(\Lambda_{ns} - \Lambda_{ls})^2 + (\Delta_{ns} + \Delta_{ls})^2]} \right), \tag{15}$$

where $\Lambda_{ns} = \mathrm{Re}\Sigma_{ns}(\varepsilon)$, $\Delta_{ns} = |\, \mathrm{Im}\Sigma_{ns}(\varepsilon)\,|$, and a_{nls} and ν_s are specified by the electronic structure of the film [see Eqs. (19) and (20) in Ref. 8]. When we employ the formalism mentioned in the previous subsection, we can self-consistently calculate the coherent potential Σ_{ns} as well as $\langle M_n^\lambda\rangle$ and $\langle(M_n^\lambda)^2\rangle$, and then the film conductivity σ by using Eqs. (13)–(15). Leaving such a detailed calculation to our future study, we adopt, in this paper, a semiphenomenological approach in discussing the temperature dependence of GMR.

We assume that magnetic layers M_1 and M_2 are separated by a nonmagnetic layer. When magnetic moments on the magnetic layers are in the antiferromagnetic (AF) and ferromagnetic (F) configurations, their conductivities are given by[8]

$$\sigma^{\mathrm{AF}} = 4c^{\mathrm{AF}}/(\Delta_\uparrow + \Delta_\downarrow), \tag{16}$$

$$\sigma^{\mathrm{F}} = c^{\mathrm{F}}[1/\Delta_\uparrow + 1/\Delta_\downarrow], \tag{17}$$

where

$$c^\eta = (e/\hbar)^2\pi\nu\, N_f^{-1} \sum_{n\in M_1} \sum_{l\in M_2} a_{nl}\tau_{nl}^\eta, \tag{18}$$

where η = AF and F, and the contribution from the interlayer scatterings between the magnetic layers is assumed to be predominant.[8] We may employ the $T = 0$ limit of Eqs. (13)–(15) because the relevant temperature is much less than the Fermi energy. Then the MR ratio is given by

$$\frac{\Delta R}{R^{\rm F}} = \frac{(R^{\rm AF} - R^{\rm F})}{R^{\rm F}} = \frac{(\sigma^{\rm F} - \sigma^{\rm AF})}{\sigma^{\rm AF}} = \frac{(a-1)^2}{4a} + \frac{b(a+1)^2}{4a(1-b)}, \quad (19)$$

with

$$a = \Delta_\uparrow/\Delta_\downarrow, \quad b = (c^{\rm F} - c^{\rm AF})/c^{\rm F}, \quad (20)$$

where R^η is the resistivity in the η (=AF, F) configuration. The first and second terms in Eq. (19) show the *short-circuit* and *valve* effects, respectively.[8] The short-circuit effect means that the total conductivity of the F states is generally larger than that of the AF state because higher-conductive ↑- (or ↓-) spin channel of the F state shortcuts the circuit, as was discussed in Refs. 3–6. On the contrary, the valve effect arises from the η dependence in τ_{nls} in Eq. (15), which plays the role of a *valve* in the transmission process between layers n and l.[8]

Now we evaluate the conductivity, taking into account the effect of spin fluctuations. The real and imaginary parts of the coherent potential in the Born approximation are given from Eqs. (6) and (7) as

$$\Lambda_s = x\left[\tilde{\varepsilon}^A - s\left(\frac{U^A}{2}\right)\langle M^A\rangle\right] + y\,\tilde{\varepsilon}^B, \quad (21)$$

$$\Delta_s = \Delta_s^r + \Delta^s, \quad (22)$$

with

$$\Delta_s^r = \pi\rho x\, y\left[\tilde{\varepsilon}^A - \tilde{\varepsilon}^B - s\left(\frac{U^A}{2}\right)\langle M^A\rangle\right]^2, \quad (23)$$

$$\Delta^s = \pi\rho\, x\left(\frac{U^A}{2}\right)^2 [\langle (M^A)^2\rangle - \langle M^A\rangle^2], \quad (24)$$

where $\tilde{\varepsilon}^\lambda = \varepsilon^\lambda + (U^\lambda/2)\langle N^\lambda\rangle$ ($\lambda = A, B$), $U^B = 0$ for a nonmagnetic B atom, and ρ is the density of states at the Fermi level. In deriving Eqs. (21)–(24), coherent potentials were assumed to be the same within the M_1 or M_2 layer, and the $2T/U_n^\lambda$ term in Eq. (12) was neglected. The first term (Δ_s^r) in Eq. (22) arises from the scatterings due to random Hartree-Fock potentials for an s-spin electron and the second term (Δ^s) comes from the effect of spin fluctuations. Equations (16)–(24) lead to

$$a = \frac{\Delta_\uparrow}{\Delta_\downarrow} = \frac{g\,(B+m)^2 + (\mu^2 - m^2)}{g\,(B-m)^2 + (\mu^2 - m^2)}, \quad (25)$$

$$b = \frac{(\Lambda_\uparrow - \Lambda_\downarrow)^2}{[(\Lambda_\uparrow - \Lambda_\downarrow)^2 + (\Delta_\uparrow + \Delta_\downarrow)^2]} \quad (26)$$

$$= \frac{m^2}{m^2 + C^{-1}\,[g\,(B^2 + m^2) + (\mu^2 - m^2)]^2} \quad (27)$$

with

$$m = \langle M^A\rangle/M_0, \quad \mu = \sqrt{\langle (M^A)^2\rangle}/M_0, \quad (28)$$

$$B = (2/U_A M_0)(\tilde{\varepsilon}_B - \tilde{\varepsilon}_A), \quad C = (2/\pi\rho\, U^A M_0)^2, \quad (29)$$

where $g = y$ and M_0 is the ground-state moment. At $T = 0$ K, a and b become

$$a_0 = a(T=0) = [(B+1)/(B-1)]^2, \quad (30)$$

$$b_0 = b(T=0) = 1/[1 + C^{-1}\, g^2\, (B^2+1)^2], \quad (31)$$

from which the coefficients B and C are expressed in terms of a_0, b_0, and g as

$$B = (\sqrt{a_0} + 1)/(\sqrt{a_0} - 1), \quad (32)$$

$$C = [b_0/(1 - b_0)]\, g^2\, (B^2 + 1)^2. \quad (33)$$

If m and μ are given, we can calculate the MR ratio with the use of Eqs. (19), (25)–(27), (32), and (33), treating a_0, b_0, and g as parameters. Physical meanings of these parameters are obvious: a_0 and b_0 are ground-state values of the asymmetry factors a and b, and g expresses the measure of the purity of the film. We may see in Eqs. (25)–(27) that $a = 1$ and $b = 0$ at $T \geq T_C$ while $a = a_0$ and $b = b_0$ at $T = 0$ K. It should be pointed out that the contribution from the so-called spin-flop process is implicitly included in Eqs. (25)–(27) through the spin-fluctuation term, which is responsible for a decrease in $\langle M^A\rangle$ when the temperature is raised. In the next section, we report numerical calculations using simple expressions for m and μ given by

$$m = \sqrt{1 - t^2}, \quad \mu = 1 \quad (t = T/T_C). \quad (34)$$

III. CALCULATED RESULTS

Calculated resistivities for AF ($R^{\rm AF}$) and F states ($R^{\rm F}$) for various a_0 with fixed values of $g = 0.1$ and $b_0 = 0$ are shown in Fig. 1(a), where they are normalized by R^C, the resistivity at the Curie temperature. When the temperature is raised, both $R^{\rm AF}$ and $R^{\rm F}$ increase. The relevant MR ratio plotted in Fig. 1(b) is shown to have a more considerable temperature dependence than the average moment, m. It is demonstrated that a film with a larger ground-state GMR has a more significantly temperature-dependent GMR. The MR ratio for the case of relatively small a_0 (= 1.5, 2) is plotted in the inset, where the GMR curves show a *quasilinear* decrease above $T/T_C \sim 0.3$ with raising the temperature. These are consistent with the results observed in many magnetic multilayers.[19–23]

We investigate the temperature dependence of the resistivity in more detail for the typical case in which $a_0 = 5$, $b_0 = 0$, and $g = 0.1$. The difference between $R^{\rm AF}$ and $R^{\rm F}$, plotted in Fig. 2(a), is proportional to $T_C - T$ at $0.3 \leq T/T_C \leq 1$. It comes from the temperature-dependent asymmetry ratio, $a = \Delta_\uparrow/\Delta_\downarrow$, shown in Fig. 2(a) where the imaginary part of the self-energy of an s-spin electron, Δ_s, is also plotted. These behaviors near the Curie temperature are easily understood from the expressions obtained from Eqs. (16), (17), (19), and (25) as

$$\Delta_s/\Delta^C \simeq 1+spm-qm^2,\ a \simeq 1+2pm+2p^2m^2,\ b \simeq rm^2, \tag{35}$$

$$R^{\mathrm{AF}}/R^C \simeq 1-(q-r)m^2,\ R^{\mathrm{F}}/R^C \simeq 1-(p^2+q)m^2, \tag{36}$$

$$\Delta R/R^{\mathrm{F}} \simeq (p^2+r)m^2,$$

with

$$p = 2gB/(gB^2+\mu^2),\ q=(1-g)/(gB^2+\mu^2), \tag{37}$$

$$r = C/(gB^2+\mu^2)^2,$$

valid for $T \lesssim T_C$ where $m^2 \propto (T_C - T)$. Figure 2(b) shows the decomposition of Δ_s to the random-potential and spin-fluctuation terms: $\Delta_s = \Delta_s^r + \Delta^s$. When the temperature is raised, $\Delta_\uparrow^r$ decreases while $\Delta_\downarrow^r$ increases since the magnetization, m, decreases. On the contrary, the spin-fluctuation term, Δ^s, monotonously increases up to the Curie temperature. The increases in R^{AF} and R^{F} are mainly due to the contribution from Δ^s.

Figures 3(a) and 3(b) show the temperature depen-

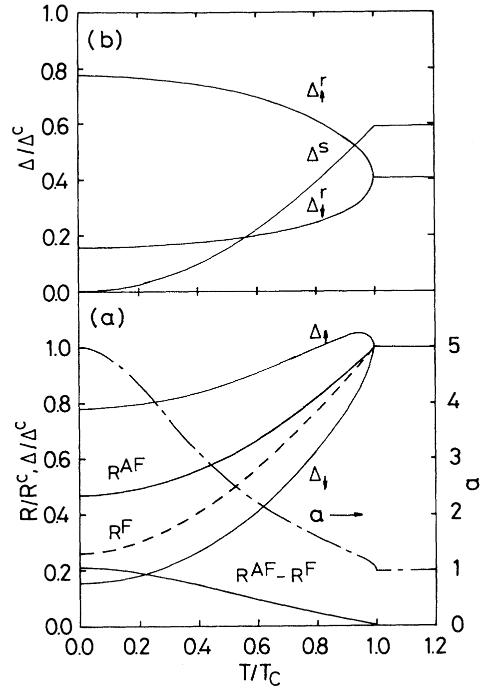

FIG. 2. (a) The temperature dependence of the resistivities of AF (R^{AF}) and F states (R^{F}), and their difference, $R^{\mathrm{AF}} - R^{\mathrm{F}}$, normalized by R^C for $a_0 = 5$, $g = 0.1$, and $b_0 = 0$. The imaginary part of the s-spin self-energy, Δ_s normalized by Δ^C, its value at $T = T_C$, and the asymmetric ratio, $a = \Delta_\uparrow/\Delta_\downarrow$ (dot-dashed curve), are also plotted. (b) The decomposition of Δ_s to $\Delta_s = \Delta_s^r + \Delta^s$.

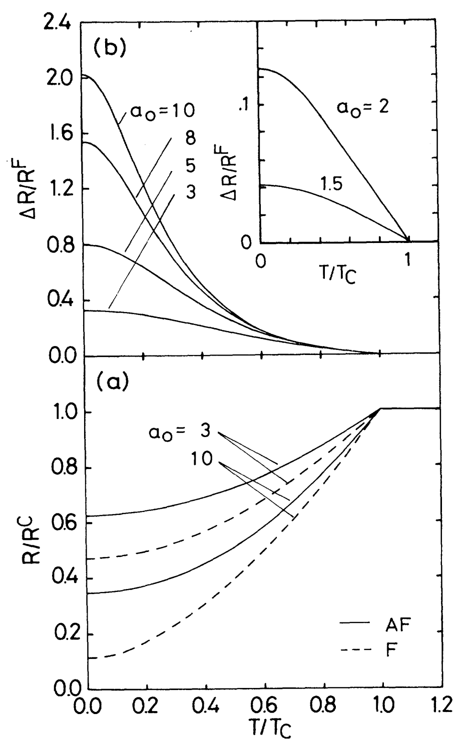

FIG. 1. (a) The temperature dependence of the resistivity R/R^C of AF (solid curves) and F states (dashed curves) normalized by R^C, the resistivity at $T = T_C$. (b) The MR ratio, $\Delta R/R^{\mathrm{F}} = (R^{\mathrm{AF}} - R^{\mathrm{F}})/R^{\mathrm{F}}$, for various a_0 with $g = 0.1$ and $b_0 = 0$, the inset showing the enlarged plot of the GMR ratio for $a_0 = 1.5$ and 2.

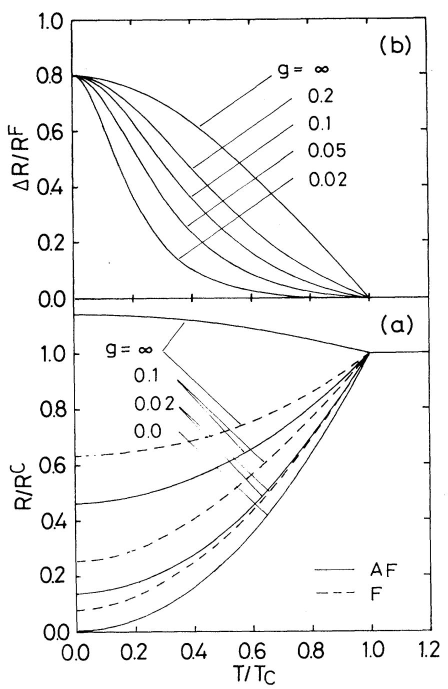

FIG. 3. (a) The temperature dependence of the resistivity and (b) the MR ratio for various g with $a_0 = 5$ and $b_0 = 0$, notations being the same as in Fig. 1.

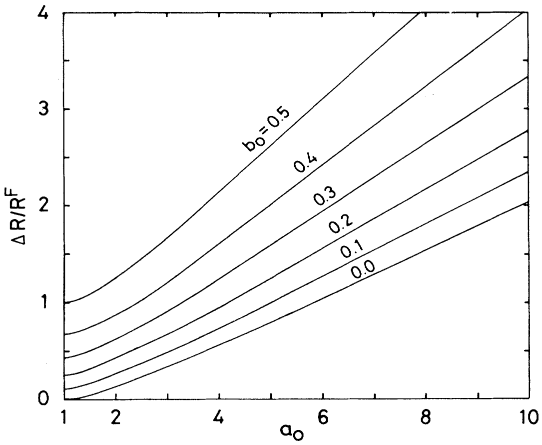

FIG. 4. The ground-state MR ratio as a function of a_0 and b_0.

dence of the resistivity and the MR ratio for various g with fixed values of $a_0 = 0.1$ and $b_0 = 0$. A film with smaller g has the lower residual resistivity as expected, and its MR ratio has more significant temperature dependence. In the limit of $g \to 0$, the contribution from the spin-fluctuation term to Δ_s is predominant, and it yields the temperature-dependent resistivity as observed in pure, bulk Fe (when the electron-phonon contribution to the resistivity is properly subtracted from the experimental data).[24] In this $g = 0$ limit, there is no difference between $R^{\rm AF}$ and $R^{\rm F}$, and the MR ratio vanishes. In the opposite limit of $g = \infty$, where there is no contribution from the spin-fluctuation term, $R^{\rm AF}$ *decreases* while $R^{\rm F}$ increases when the temperature is raised [see $\Delta_\uparrow^r$ and $\Delta_\downarrow^r$ in Fig. 2(b)]. According to experimental data, however, both $R^{\rm AF}$ and $R^{\rm F}$ increase with raising the temperature.[20–22] Our calculation shows that the spin-fluctuation contribution, Δ^s, is crucial to account for the observed temperature dependence of the resistivity ($R^{\rm AF}, R^{\rm F}$) and GMR, while the random-potential contribution, Δ_s^r, is indispensable to yield the GMR.

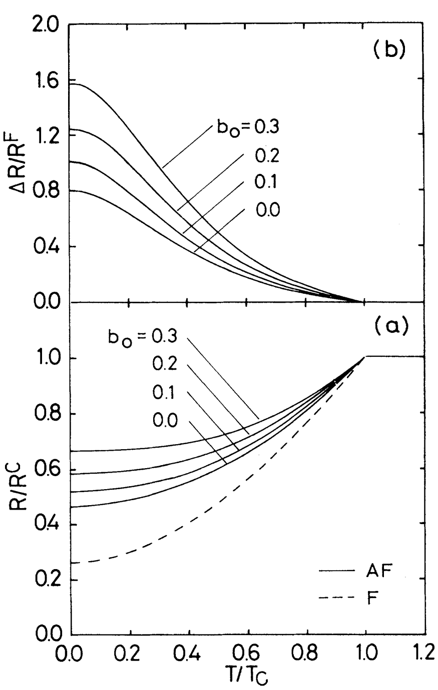

FIG. 5. (a) The temperature dependence of the resistivity and (b) the MR ratio, for various b_0 with $a_0 = 5$ and $g = 0.1$, notations being the same as in Fig. 1.

So far we have studied the short-circuit effect given by the first term of Eq. (19). In order to investigate the valve effect expressed by its second term, we made calculations with finite b_0 values. Figure 4 shows the ground-state MR ratio as a function of a_0 and b_0. The valve effect works to enhance the GMR. The calculated temperature dependence of resistivity and the MR ratio are plotted in Figs. 5(a) and 5(b), where the b_0 value is changed but with $a_0 = 5$ and $g = 0.1$. As b_0 is increased, the ground-state MR ratio is increased and its temperature dependence becomes significant.

IV. CONCLUSION AND DISCUSSION

We have discussed the temperature dependence of the resistivity and the GMR in magnetic multilayers. We have elucidated the mechanism of their temperature dependence, showing that spin fluctuations play primary roles at finite temperatures. Numerical calculations using the simple model demonstrate that the temperature dependence of the calculated GMR is much larger than that of the average magnetization. The temperature variation of the MR ratio is shown to become more considerable in a multilayer with the larger ground-state MR ratio (larger a_0 and/or b_0) and/or better purity (smaller g). Our calculation accounts for, at least qualitatively, the feature observed in many transition-metal multilayers.[19–23]

Because our spin-fluctuation theory[11] employs the static approximation in evaluating the functional integral, we cannot draw any definite conclusion on the temperature dependence of the resistivity (and the GMR) at low temperatures, where dynamical effect of spin fluctuations yields the T^2 contribution.[25] Although we have taken into account only the scatterings due to random potentials and spin fluctuations, there are many other mechanisms to be responsible for the resistivity. For example, the intraband and interband electron-phonon scatterings are well known to yield the T^5 and T^3 resistivities, respectively,[26] at low temperatures, where the electron-electron Baber scatterings lead to a T^2 resistivity.[27] Nevertheless, we expect that the mechanisms examined in the present paper are most important in transition metals like Fe,[24] and that our calculation provides us with an *overall* picture on the temperature-dependent GMR of magnetic multilayers.

We note in Eqs. (19), (20), and (25)–(27) that the temperature dependence of GMR depends on those of m and μ. Although our model calculation adopted a simple expression given by Eq. (34), the MR ratio in real systems may show the temperature dependence different

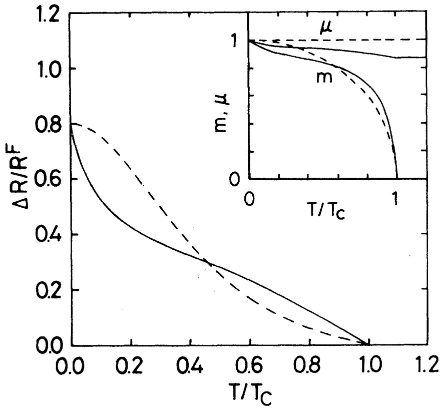

FIG. 6. The calculated MR ratio when the average (m) and rms (μ) moments are given as shown in the inset where the solid curves denote the results of Ref. 28 and the dashed curves those given by Eq. (34) (see text).

from our results, if the temperature dependence of their m and μ are different from that in our model. Quite recently, Hasegawa[28] calculated the temperature dependence of the average of the magnetic moment, $\langle M_n \rangle$, and its root-mean-square value, $\sqrt{\langle M_n^2 \rangle}$, on layer n of a realistic Fe/Cr multilayer using the finite-temperature band theory.[11,17] The average moment on the central Fe layer approximately follows that in bulk Fe. On the other hand, the moment on the interface Fe layer shows a peculiar temperature dependence, which is plotted in the inset of Fig. 6. This is expected to arise from weaker magnetic couplings between Fe-Cr layers than those between Fe-Fe layers. Similar results are obtained in Co/Cr multilayers.[17] If we use these temperature-dependent m and μ shown in the inset,[28] we obtain the GMR plotted in Fig. 6, which shows a rather peculiar temperature dependence, in particular at low temperatures. This suggests that the temperature dependence of $\Delta R/R^{\mathrm{F}}$ in real systems shows a variety depending on the temperature dependence of the $\langle M_n \rangle$ and $\langle M_n^2 \rangle$. It might be possible that some information on the temperature dependence of $\langle M_n \rangle$ and $\langle M_n^2 \rangle$ is obtainable from the observed temperature-dependent MR ratio. A detailed calculation of the finite-temperature GMR using a realistic tight-binding model is now in progress, and will be reported separately.

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