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The influence of the bound d state on the magnetoresistance in magnetic multilayers

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Abstract. We discuss the giant magnetoresistance (GMR) in magnetic superlattices with an s–d scattering model on the assumption that d states are bound in magnetic layers. The GMR is calculated by using the quantum Boltzmann equation using Kronig–Penney-type potentials. Spin-dependent interfacial scattering depends on the number of scatterers, the height of the scattering potentials and the amplitude of the wave function of the d state at interfaces, while spin-dependent bulk scattering is attributed to the spin-dependent density of states (DOS) of d states. Our model agrees well with the measured GMR in Co/Cu superlattices with artificially mixed interfaces, when we assume that the minority-spin d states are strongly bound in Co layers. Therefore, the spin-dependent scattering in Co/Cu superlattices is attributed to the spin-dependent DOS of the d states in the Co layers.

1. Introduction

The giant magnetoresistance (GMR) exhibited in many magnetic superlattices and granular alloys arises from the spin-dependent scattering of conduction electrons [1–5]. It has been pointed out that the interfacial scattering of conduction electrons plays an important role in the spin-dependent scattering [1]. Camley and Barnaś [6, 7] have proposed a semi-classical model for the GMR by extending the Fuchs–Sondheimer model [8]. In their model, the GMR is described with many phenomenological parameters for conduction electrons such as spin-dependent mean free paths in layers, reflection, transmission and diffuse scattering coefficients at interfaces. Quantum mechanical models by Hood and Falicov [9] and Visscher [10] relate these parameters to superlattice potentials and the potentials of bulk and interfacial scatterers. Since these theories are based on a single-band free electron model, spin-dependent scattering is attributed to spin-dependent potentials. For Fe/Cr superlattices, the spin-dependent potentials due to interfacial roughness have been indicated by microscopic theories [11, 12]. In fact, it has been reported that the magnetoresistance (MR) ratio in Fe/Cr is enhanced drastically by the interfacial roughness [13, 14]. Thus, the origin of the GMR in Fe/Cr is the spin-dependent scattering at interfaces.

For M/Cu (M = Co, NiFe) systems, however, no one has reported the enhancement of the GMR owing to the interfacial roughness [15–19]. In our previous paper [16], we have reported that the temperature dependence of the GMR is not influenced by the interfacial roughness, and the residual MR ratio decreases with increasing interfacial roughness. This suggests that the spin-dependent bulk scattering is important for the occurrence of the GMR. On the other hand, Parkin [20] has reported that the GMR is enhanced by very thin Co layers inserted at interfaces of NiFe/Cu multilayers. He claimed that the *interfacial state* as well as roughness plays an important role in the GMR. These experimental results require a theory

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to take account of both the spin-dependent bulk scattering and the interfacial electronic states. In discussing the GMR in M/Cu, however, the above theories are not appropriate, since they do not include the information about the band structure being very different from that of Fe/Cr. If we take account of the electronic band structure, the formalism for the bulk and interfacial scattering will be modified. For M/Cu systems, Edwards *et al* [21] claimed in their resistor network theory that the spin-dependent s–d scattering in bulk is the dominant process giving rise to the GMR. Xing *et al* [22] also insist on the importance of a spin-dependent density of states (DOS) of d bands in magnetic layers. In these models, the GMR is attributed to the scattering of s electrons to unfilled d bands which have spin-dependent DOS. These models semi-quantitatively agree with the layer thickness dependence of GMR in M/Cu systems. However, the influence of the interfacial state has not been treated explicitly, since the Fermi surfaces of s and d bands are treated as simple spheres. Recently, Schep *et al* [23] calculated the GMR in a method based on the full electronic structure. They reported the importance of the s–d hybridization for the origin of the GMR in the current-perpendicular-to-the-plane (CPP) geometry. However, their theory does not satisfactorily explain the considerable MR observed in the current-in-the-plane (CIP) geometry. They suggest that some additional scattering mechanism is necessary for explaining the CIP MR.

In this study, we have extended the resistor network theory to include the interfacial state explicitly, on the basis of the method developed by Visscher [10]. The simple expression we have deduced neglecting the s–d hybridization agrees well with our experimental results for the GMR observed in CIP geometry for the Co/Cu superlattices with artificially mixed interfaces.

2. The model

We confine our discussions to low temperatures, neglect magnon and phonon scattering and assume that the effective mean free path of the conduction electrons is much larger than the superlattice period. This assumption limits our discussion to scattering processes that do not cause spin mixing, so that the current is carried separately by up- and down-spin conduction electrons (two-current model) as [24]

$$\rho = \frac{\rho_\uparrow \rho_\downarrow}{\rho_\uparrow + \rho_\downarrow} \tag{1}$$

where $\rho_\uparrow$ and $\rho_\downarrow$ are the resistivity of up- and down-spin channels, respectively, and ρ is the total resistivity. We assume that the current is carried only by s electrons and neglect the current carried by d electrons with heavy effective mass. The non-zero residual resistivity is due to the scattering of s electrons by impurities or defects at r_{SC} with a potential of the form

$$V(r) = V_0\,\delta(r - r_{SC}). \tag{2}$$

In the transition metals, the s electrons can be scattered into holes in the s band (s–s scattering) or d band (s–d scattering) when they are scattered either at interfaces or inside layers. The scattering rate of s–s and s–d scattering will be determined by the scattering amplitude for the transition between an initial and a final state and the availability of states into which the electrons can be scattered. Since only electrons near the Fermi surface contribute to the resistivity, these scattering rates can be written in terms of Fermi's golden rule as

$$P_{ij}^{\sigma} = \frac{2\pi}{\hbar}|\langle \Phi_j^{\sigma}|V(r)|\Psi_i^{\sigma}\rangle|^2 N(E_F) \tag{3}$$

where Ψ_i^σ and Φ_f^σ are an initial s and a final s or d state for the σ spin, respectively, and $N(E_F)$ is the DOS of the final state at the Fermi energy E_F. The total scattering rate is the summation of the s–s and s–d scattering rates. Therefore, the properties of s and d bands near the Fermi surface play an important role in the GMR. We treat the s and d bands with the effective-mass approximation. Since the CIP MR due to s–d hybridization is small [23], we neglect it and assume that the bulk of each metal constituting the superlattice has two free electron bands with different effective masses. One corresponds to an s band, and another corresponds to a d band. The effective masses and potentials for the bands may depend on spin directions for the ferromagnetic metals. When two metals are layered alternately, the superlattice states will be formed owing to the potential modulation. For systems such as M/Cu, s electrons see small potential differences at interfaces between magnetic and nonmagnetic layers, since the bottoms of the s bands of these metals are well aligned on the scale of the Fermi energy [25, 26]. Owing to the superlattice state due to the superlattice potential, the anisotropy in the Fermi wave number will arise. Although this anisotropy affects the CPP MR, its influence on the CIP MR can be negligible [27]. This means that the s states can be treated as simple plane waves when the potential differences are small. For simplicity, we assume that s electrons do not see any potential differences when crossing the interfaces between magnetic and non-magnetic layers, since we are now interested in the CIP MR. On the other hand, d states see a large potential difference at interfaces, since there is an exchange interaction in ferromagnetic layers. Assuming the superlattice to be infinite in the z direction, the superlattice potentials are modelled by Kronig–Penney-type potentials. We assume that the potentials for the d states are constant within each layer but depend on spin and material. For systems such as M/Cu, the d band of Cu and the majority-spin d band of the ferromagnetic layers lie well below the Fermi energy, so they do not contribute to the resistivity [25, 26]. Therefore, we take account of only the minority-spin d band of ferromagnetic layers. The potential for the d band is composed of a periodic array of barriers of thickness a with height U_b and wells of thickness b with depth U_w as shown in figure 1. For a ferromagnetic (F) configuration of the magnetization of ferromagnetic layers, the potential for down-spin electrons has $a = t_{NM}$ and $b = t_M$, where t_{NM} is the thickness of nonmagnetic layers and t_M of magnetic layers; for up-spin electrons, $b = 0$, i.e., there are no wells. For antiferromagnetic (AF) configurations, the potentials for up- and down-spin electrons are the same and are shifted in space relative to one another; $a = 2t_{NM} + t_M$ and $b = t_M$. The potential for the s band is at U_s and is constant in a superlattice as mentioned above. For these step-function potentials, we calculate the scattering matrix elements from the exact quantum-mechanical wave functions.

3. The superlattice state for a d band

Since the potential for s electrons is constant in the superlattice, s electrons travel in the superlattice as ordinary plane waves. Thus, the Fermi surface of the s band is spherical. However, the Fermi surface of the d band is no longer spherical due to the superlattice potential. In order to calculate the resistivity, we must know the shape of the Fermi surface of the d band. The wave functions of the d state $\Psi(r)$ have the Bloch form

$$\Psi(r) = \psi(z)\exp(ik_x + ik_y) \tag{4}$$

where $\psi(z)$ satisfies the Bloch condition for some k_z

$$\psi(z + D) = \exp(ik_z D)\psi(z) \tag{5}$$

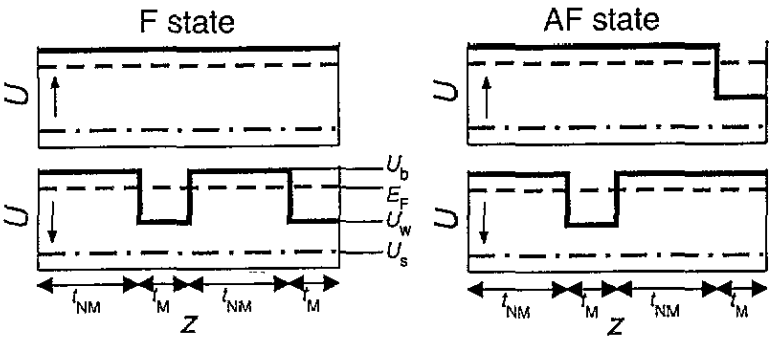

Figure 1. A schematic diagram of the superlattice potentials for ferromagnetic (F state) and antiferromagnetic (AF state) configurations and for electrons with up (↑) and down (↓) spin. The lines indicate the potentials for d states (——), the potentials for s states (— · —) and the Fermi energy (– – –).

where $D = a + b$. The Brillouin zone is infinite in the x and y directions, and $2\pi/D$ wide in z (k_z) direction. Here, we normalize the wave function Ψ as

$$\Omega^{-1} \int |\Psi(r)|^2 \, \mathrm{d}^3 r = 1 \tag{6}$$

where Ω is the volume of a normalization box. The solution of $\psi(z)$ is a linear combination of $\exp(\pm \mathrm{i} k_L z)$, where k_L on the Fermi surface is defined in the well as

$$k_L = \left[\frac{2m_\mathrm{d}}{\hbar^2}(E_\mathrm{F} - U_\mathrm{w}) - {k_r}^2 \right]^{1/2} \equiv K \tag{7}$$

and in the barrier as

$$k_L = \left[\frac{2m_\mathrm{d}}{\hbar^2}(E_\mathrm{F} - U_\mathrm{b}) - {k_r}^2 \right]^{1/2} \equiv \mathrm{i}Q \tag{8}$$

where $k_r = ({k_x}^2 + {k_y}^2)^{1/2}$ and m_d is the effective mass of the d states. From (5) and the boundary conditions where ψ and $\mathrm{d}\psi/\mathrm{d}z$ are continuous, we obtain

$$\frac{\mathrm{e}^{Qb}}{2} \left[\frac{Q^2 - K^2}{2QK} \sin(Ka) + \cos(Ka) \right] = \cos(k_z D). \tag{9}$$

In this study, we consider a case in which the barrier height is large enough to satisfy $\mathrm{e}^{Qb} \gg 1$. In this case, the minority-spin d states are confined to the Co layers and have eigenstates quantized to the thickness of Co layers. Therefore, the Fermi surface becomes a set of cylinders (subbands) parallel to the z (or k_z) direction.

4. Calculation of the resistivity

Resistivities are calculated by the relaxation time approximation of the Boltzmann equation. Since we neglect the shift of the Fermi surface for the d states, the relaxation time $\tau_\sigma(k)$ ($\sigma = \uparrow, \downarrow$) is given as [22, 24]

$$\frac{1}{\tau_\sigma(k)} = \frac{1}{\tau_\sigma^{\mathrm{ss}}(k)} + \frac{1}{\tau_\sigma^{\mathrm{sd}}(k)} \tag{10}$$

$$\frac{1}{\tau_\sigma^{\mathrm{ss}}(k)} = \int \left(1 - \frac{k \cdot k'}{k^2} \right) P(k'\mathrm{s}\sigma, k\mathrm{s}\sigma) \, \mathrm{d}^3 k' \tag{11}$$

$$\frac{1}{\tau_\sigma^{sd}(k)} = \int P(k'd\sigma, ks\sigma)\, d^3k' \tag{12}$$

where $P(k'i, kj)$ is the scattering rate between states $(k'i)$ and (kj), i and j stand for s↑, s↓, d↑ and d↓. The first term in (10) comes from s–s scattering, and the second from s–d scattering. For simplicity, we assume here that the bulk scatterers with the potential strength $V_0 = V_B$ distribute in the superlattices with uniform volume density of η_B. In addition, we assume that the interfacial scatterers with the potential strength $V_0 = V_I$ distribute within the region of $z = z_i \pm \varepsilon$ with the uniform volume density of η_I, where z_i is the position of the ith interface. Here, we do not assume any spin dependence in V_B and V_I, as in [21] and [22].

4.1. s–s scattering

Since the wave function of the s state is a plane wave, it is easy to calculate the rate of s–s scattering. Using Fermi's golden rule, the scattering rate is given as [10]

$$P(k's\sigma, ks\sigma)d^3k' = \frac{2\pi}{\hbar}\left(V_B{}^2\eta_B + n_I^{ss}\frac{2\varepsilon}{D}V_I{}^2\eta_I\right) dN(E_F) \tag{13}$$

where n_I^{ss} is the number of interfaces in a period of superlattice potential and $dN(E_F)$ denotes the number of states per unit volume and unit energy at E_F in the volume element d^3k' near k'. Integrating (13), we obtain the relaxation time due to s–s scattering as

$$\frac{1}{\tau_\sigma^{ss}} = \frac{2\pi}{\hbar} g_s(E_F)\left(V_B{}^2\eta_B + n_I^{ss}\frac{2\varepsilon}{D}V_I{}^2\eta_I\right) \tag{14}$$

$$g_s(E_F) = \frac{m_s k_F^s}{2\pi^2\hbar^2} \tag{15}$$

where m_s and k_F^s are the effective mass and the Fermi wave number of s electrons. The first term in the parentheses of (14) comes from the bulk scattering, while the second term is due to the interfacial scattering.

4.2. s–d scattering

The scattering rate of s–d scattering is written as

$$P(k'd\sigma, ks\sigma)\, d^3k' = \frac{2\pi}{\hbar}\left(V_B{}^2\eta_B + \sum_i \frac{2\varepsilon|\psi_{k'}(z_i)|^2}{D}V_I{}^2\eta_I\right) dN(E_F) \tag{16}$$

when $\varepsilon \ll 1$. The summation in (16) is taken over the interfaces in a period of the superlattice potential. In the case where $Qb \gg 1$, d states are confined in the well layers, so the amplitude of $\psi_{k'}(z)$ has a non-zero value in the well layers and near the interfaces between the well and the barrier layers. Thus, it is enough to take the summation in (16) over the well–barrier interfaces. From the analogy with the simple quantum well problem, $|\psi_{k'}(z_i)|^2$ has the same value at the interfaces as at both edges of the well layers. Hence, (16) is written with the number of the well–barrier interfaces in a period n_I^{sd} as

$$P(k'd\sigma, ks\sigma)d^3k' = \frac{2\pi}{\hbar}\left(V_B{}^2\eta_B + n_I^{sd}\frac{2\varepsilon|\psi_{k'}(z_I)|^2}{D}V_I{}^2\eta_I\right) dN(E_F) \tag{17}$$

where z_I is the position of one of the interfaces causing non-zero s–d scattering. As mentioned in the previous section, allowed-wave-number vectors k' are discrete. In each

subband, the value of k_r' and the wave function are constant, since $Qb \gg 1$. Thus, we obtain

$$dN(E_F) = \frac{m_d}{(2\pi\hbar)^2}\,dk_z' \tag{18}$$

for each subband, where dk_z' is a length increment along the Fermi surface in the z direction. Integrating (17) using (18), the relaxation time due to s–d scattering is given with the k_r' in the jth subband $k_r'^j$ as

$$\frac{1}{\tau_\sigma^{sd}} = \frac{2\pi}{\hbar} g_d(E_F)\frac{a}{D}\left(V_B{}^2\eta_B + n_I^{sd}\frac{2\varepsilon\phi}{D}V_I{}^2\eta_I\right) \tag{19}$$

$$\phi = \frac{1}{n_d}\sum_{j=1}^{n_d}\left|\psi_{k_r'^j}(z_I)\right|^2 \tag{20}$$

$$g_d(E_F) = \frac{m_d n_d}{2\pi\hbar^2 a} \tag{21}$$

where n_d is the number of d subbands. The first term in the parentheses of (19) comes from the bulk scattering, while the second term is due to the interfacial scattering. Since the relaxation time due to both s–s and s–d scattering (equations (14) and (19)) is independent of k, the resistivity is written with the density of s electrons N_s as

$$\rho_\sigma = \frac{2\pi m_s}{\hbar N_s e^2}\left\{g_s(E_F)\left(V_B{}^2\eta_B + n_I^{ss}\frac{2\varepsilon}{D}V_I{}^2\eta_I\right) + g_d(E_F)\frac{a}{D}\left(V_B{}^2\eta_B + n_I^{sd}\frac{2\varepsilon\phi}{D}V_I{}^2\eta_I\right)\right\}. \tag{22}$$

Equation (15) indicates that $g_s(E_F)$ is the DOS of the s states per spin at the Fermi energy for the bulk. On the other hand, $g_d(E_F)$ depends on n_d due to the quantum size effect. When the barrier height U_b is large enough, n_d is approximated by the maximum integer satisfying $n_d \leqslant (a/\pi)\left[2m_d(E_F - U_w)/\hbar^2\right]^{1/2}$. If we neglect the discreteness in n_d, $g_d(E_F)$ is identical with the DOS of the d states at the Fermi energy for the bulk ferromagnetic metals. In this case, (22) has the same form as the resistor network theory except for the terms due to the interfacial scattering. The interfacial scattering depends on the number of the scattering centres ($2\varepsilon\eta_I$), the strength of the scattering potential (V_I) and the amplitude of the wave function of the d state at interfaces ($|\psi_k(z_I)|^2$). The first one corresponds to the magnitude of the interfacial roughness, and the last one corresponds to interfacial electronic state. Therefore, our model includes not only the interfacial roughness but also the interfacial electronic states.

5. GMR

In this section, we calculate the GMR using (22). The parameters dependent on the spin and magnetic configuration are D, n_I^{ss}, n_I^{sd} and ϕ. Defining the superlattice period as $\lambda = t_M + t_{NM}$, the value of D for the F state is λ and that for the AF state is 2λ. According to this, the value of n_I^{ss} for the F state is 2, and that for the AF state is 4. The values of n_I^{sd} for the down-spin electrons for the F state and for the up- and down-spin electrons for the AF state are 2, while $n_I^{sd} = 0$ for up-spin electrons for the F state. Since the normalization condition depends on the superlattice period (see (6)), ϕ for the AF state is two times larger than that for the F state. Thus, $\phi = 2\phi_F$ for the AF state, where ϕ_F is the ϕ for the F state. The number of the subbands n_d is constant, unless the thickness of the magnetic layers changes. With these parameters, we obtain the expressions for the spin-dependent

resistivities for the F state ρ_σ^{F} as

$$\rho_\uparrow^{\mathrm{F}} = \frac{2\pi m_{\mathrm{s}}}{\hbar N_{\mathrm{s}} e^2}\left[g_{\mathrm{s}}(E_{\mathrm{F}})\left(V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} + \frac{4\varepsilon}{\lambda}V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}\right)\right] \tag{23}$$

$$\rho_\downarrow^{\mathrm{F}} = \frac{2\pi m_{\mathrm{s}}}{\hbar N_{\mathrm{s}} e^2}\left[g_{\mathrm{s}}(E_{\mathrm{F}})\left(V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} + \frac{4\varepsilon}{\lambda}V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}\right) + g_{\mathrm{d}}(E_{\mathrm{F}})\frac{t_{\mathrm{M}}}{\lambda}\left(V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} + \frac{4\varepsilon\phi_{\mathrm{F}}}{\lambda}V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}\right)\right] \tag{24}$$

and the resistivities for the AF state $\rho_\sigma^{\mathrm{AF}}$ as

$$\rho_\uparrow^{\mathrm{AF}} = \rho_\downarrow^{\mathrm{AF}} = \frac{2\pi m_{\mathrm{s}}}{\hbar N_{\mathrm{s}} e^2}\left[g_{\mathrm{s}}(E_{\mathrm{F}})\left(V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} + \frac{4\varepsilon}{\lambda}V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}\right) + g_{\mathrm{d}}(E_{\mathrm{F}})\frac{t_{\mathrm{M}}}{2\lambda}\left(V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} + \frac{4\varepsilon\phi_{\mathrm{F}}}{\lambda}V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}\right)\right]. \tag{25}$$

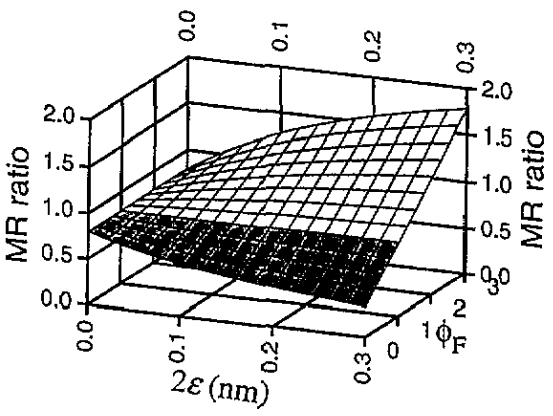

Figure 2. The calculated MR ratio with respect to 2ε and ϕ_{F}. The parameters are assumed to be $t_{\mathrm{M}} = 1$ nm, $t_{\mathrm{NM}} = 2$ nm, $g_{\mathrm{d}}(E_{\mathrm{F}})/g_{\mathrm{s}}(E_{\mathrm{F}}) = 12$ and $V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}/V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} = 5$. The regions of $\phi_{\mathrm{F}} < 1$ and $\phi_{\mathrm{F}} > 1$ are indicated by the different tones.

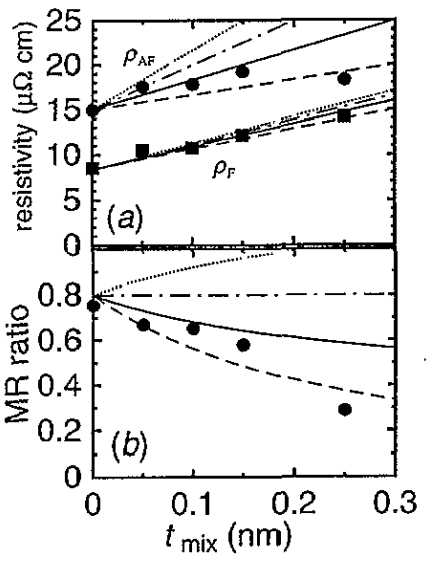

Figure 3. The measured (symbols) and calculated (lines) resistivities (*a*) and MR ratios (*b*). Measurement was performed at 5 K on $[\mathrm{Co}(1.0\ \mathrm{nm})/\mathrm{Cu}(2.2\ \mathrm{nm})]_{16}$ with artificially mixed interfacial regions of thickness t_{mix}. In (*a*), the measured ρ_{AF} (●) and ρ_{F} (■) correspond to the resistivities for antiferromagnetic and ferromagnetic configurations, respectively. Data are taken from [16]. Parameters used in the calculation are $t_{\mathrm{M}} = 1$ nm, $t_{\mathrm{NM}} = 2.2$ nm, $g_{\mathrm{d}}(E_{\mathrm{F}})/g_{\mathrm{s}}(E_{\mathrm{F}}) = 12$ and $V_{\mathrm{I}}{}^2\eta_{\mathrm{I}}/V_{\mathrm{B}}{}^2\eta_{\mathrm{B}} = 5$. The lines indicate the results for $\phi_{\mathrm{F}} = 0$ (– – –), $\phi_{\mathrm{F}} = 0.5$ (——), $\phi_{\mathrm{F}} = 1.0$ (— · —) and $\phi_{\mathrm{F}} = 1.5$ (· · · · · ·).

Although the GMR can be calculated numerically assuming the appropriate superlattice potentials, here we discuss the GMR phenomenologically, treating $g_{\mathrm{d}}(E_{\mathrm{F}})$ and ϕ_{F} as independent parameters for better transparency of the physics. Figure 2 shows the calculated MR ratio with respect to 2ε and ϕ_{F}. We assume here that $t_{\mathrm{M}} = 1$ nm, $t_{\mathrm{NM}} = 2$ nm,

$g_d(E_F)/g_s(E_F) = 12$ and $V_I^2\eta_I/V_B^2\eta_B = 5$. These values are suitable for explaining the experimental results for Co/Cu superlattices, as we mention in the next section. A considerable MR ratio is obtained even without the interfacial scattering ($2\varepsilon = 0$ nm) because of the strong spin dependence in the DOS of the d states. The MR ratio decreases with increasing 2ε for $\phi_F \leqslant 1$, while it increases with increasing 2ε for $\phi_F \geqslant 1$. This tendency remains unchanged if we use different values for the layer thicknesses, i.e. the ratios of the DOS and the interfacial to bulk scattering remains unchanged. Thus, the relationship between the MR ratio and the interfacial roughness strongly depends on ϕ_F. In other words, the GMR in superlattices with interfacial roughness strongly depends on the amplitude of the wave function of the d state at interfaces. In the next section, we estimate the magnitude of ϕ_F for Co/Cu superlattices by comparing our calculation with experiments.

6. Interfacial scattering in Co/Cu

We have reported the structure and the GMR properties of Co/Cu superlattices with artificially mixed interfaces [15, 16, 28]. The interfaces in Co/Cu have been modified by codeposition. The Co and Cu atoms are mixing randomly in the interfacial regions, which increases with increasing nominal thickness of the codeposited region t_{mix}.

The resistivity of these samples decreases from the initial value at zero field with increasing magnetic field, and saturates at the value of ρ_F in a field larger than the saturation field. After saturation, the resistivity has a peak near the coercive field. Since the value of the initial resistivity is larger than that of the peak one, the AF alignment of the magnetization of the Co layers is more perfect in the initial state than that in the field where the resistivity has the peak. The magnetization of all samples is zero in the initial state, and the relationship between the resistivity and the square of the magnetization is linear at low temperatures. From the discussion in [16], this indicates that the magnetic configuration of our samples is very close to the perfect AF alignment in the initial state. Therefore we denote the resistivity in the initial states as ρ_{AF} here. Furthermore, the linear dependence of the resistivity on the square of the magnetization precludes the strong spin dependence of the scattering potential for both bulk and interfacial scattering.

Figure 3 shows the t_{mix} dependence of the GMR measured at 5 K for [Co(1.0 nm)/Cu(2.2 nm)]$_{16}$. Data are taken from [16]. Both ρ_{AF} and ρ_F increase with increasing t_{mix}, while the MR ratio decreases. It is clear that the interfacial scattering is crucial but less spin dependent.

Despite the large differences in the resistivities at low temperature, no significant differences in the temperature dependence of the GMR between the samples with different t_{mix} have been observed. For all samples, ρ_F closely approximates a T^2 power law, while $\rho_{AF} - \rho_F$ changes linearly with $T^{3/2}$. The difference between the temperature coefficients of the samples with $t_{mix} \leqslant 0.15$ nm is small. This indicates that the interfacial scattering can be attributed to the impurity or defect scattering.

Although the mean free path roughly estimated for the maximum resistivity in figure 3 ($\approx 20\ \mu\Omega$ cm) is about 4.0 nm, Edwards *et al* [21] indicated that the limit of long mean free path is already reached rapidly for a mean free path comparable with the superlattice period. Therefore, it is appropriate to adopt the present model for interpreting the t_{mix} dependence of the GMR at 5 K. We can estimate the value of $g_d(E_F)/g_s(E_F)$ at about 12 from the measured MR ratio of the sample of $t_{mix} = 0$ nm, if we neglect the interfacial scattering (assuming $\varepsilon = 0$ nm in equations (23)–(25)). This value of $g_d(E_F)/g_s(E_F)$ is close to the ratio of the calculated DOS at E_F in the majority- and minority-spin bands in bulk Co [25, 26]. The calculated resistivities and the MR ratio for the parameters of $t_M = 1$ nm,

$t_{NM} = 2.2$ nm, $g_d(E_F)/g_s(E_F) = 12$ and $V_I^2\eta_I/V_B^2\eta_B = 5$ are also indicated by lines in figure 3, where we assume that $t_{mix} = 2\varepsilon$. The calculated results are normalized at $t_{mix} = 0$ so as to make the calculated ρ_F equal to the measured one. It is clear that the behaviour of ρ_F is insensitive to ϕ_F. The agreement between measured and calculated ρ_F is good when $V_I^2\eta_I/V_B^2\eta_B = 5.0 \pm 1.0$. It is reasonable that the density of the scattering centres is larger, or the scattering potential is stronger at interfaces than in the bulk ($\eta_I > \eta_B$ or $V_I > V_B$). On the other hand, ρ_{AF} is sensitive to ϕ_F. Comparing the measured and the calculated results, $\phi_F \simeq 0.5$ is suitable for explaining the behaviour of ρ_{AF} and the MR ratio for the samples for which $t_{mix} \leqslant 0.15$ nm. The approximation $\varepsilon \ll 1$ in (16) will no longer hold when t_{mix} becomes comparable with the period of $|\psi_k(z)|^2$. This affects ρ_{AF} much more than it does ρ_F, and may cause the deviation between the measured and calculated ρ_{AF} at $t_{mix} = 0.25$ nm. Detailed numerical calculation will be necessary to study the GMR for the samples with such large roughness. The value of $\phi_F < 1$ indicates that amplitude of the d states is attenuated at interfaces, since the average value of $|\psi_k(z)|^2$ is about λ/t_M (> 1) in the magnetic layers. Thus, d states are strongly bound in magnetic layers.

Our model agrees well with the experimental results. However, the contribution of the interfacial scattering to the resistivity of the sample with $t_{mix} = 0$ nm remains unclear, although we neglect it in the above discussion. In fact, we have confirmed the existence of a small amount of interfacial mixing for the sample with $t_{mix} = 0$ nm [28]. The influence of the interfacial mixing can also be discussed by the layer thickness dependence of the resistivities. Thus, we focus our attention on the dependence of ρ_F (which is not influenced by the interlayer coupling between adjacent Co layers) on the thickness of the Co and Cu layers for samples without intentionally mixed interfaces. As a result, ρ_F increases with increasing Co layer thickness, while it decreases with increasing Cu layer thickness. This indicates clearly that the resistivity of Co layers is larger than that of not only Cu layers but also the interfaces. This layer thickness dependence is reproduced only when ε is very small in equations (23) and (24). Therefore, our assumption of $\varepsilon = 0$ for the sample with $t_{mix} = 0$ nm is appropriate.

If superlattices have slightly rougher interfaces than those of our samples with $t_{mix} = 0$ nm, the interfacial scattering will influence the GMR significantly depending on the interfacial electronic state. The interfacial electronic state should be sensitive to the combination of metals. Parkin [20] has reported that the MR ratio is enhanced by inserting very thin Co layers at the interfaces between Cu and NiFe layers. This can be interpreted in our model as the change in ϕ_F depending on materials. The virtual bound state indicated by Inoue and Maekawa [29] may also affect the magnitude of the GMR. However, the bulk s–d scattering will be the most important process giving rise to the GMR in M/Cu superlattices, since a very large MR ratio is observed, at least for our sample whose interfacial scattering is very small.

Consequently, the GMR in Co/Cu superlattices mainly comes from the spin-dependent bulk scattering due to spin-dependent DOS in the Co layers. The interfacial scattering due to interfacial roughness is less spin dependent than the bulk scattering. This is understood in terms of the d states bound in the Co layers.

7. Conclusions

We have extended the resistor network theory to include interfacial scattering explicitly and discussed the role of bound d state in the CIP MR. The interfacial states are described in our model with the number of interfacial scattering centres ($2\varepsilon\eta_I$), the height of the scattering potentials (V_I) and the amplitude of the wave function of the d states ($|\psi_k(z_I)|^2$). The

first is concerned with the interfacial roughness, and the last corresponds to the interfacial electronic state. Our model is consistent with the experiments for Co/Cu superlattices with artificially mixed interfaces, when we assume that the minority spin d states are strongly bound in Co layers. Therefore, the GMR in Co/Cu superlattices mainly comes from the spin-dependent s–d scattering in the Co layers, and the interfacial scattering is less spin-dependent than bulk scattering. This is because of the attenuation of the wave function of d states at interfaces.

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