Conductance and exchange coupling of two ferromagnets separated by a tunneling barrier

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A theory is given for three closely related effects involving a nonmagnetic electron-tunneling barrier separating two ferromagnetic conductors. The first is Julliere's magnetic valve effect, in which the tunnel conductance depends on the angle θ between the moments of the two ferromagnets. One finds that discontinuous change of the potential at the electrode-barrier interface diminishes the spin-polarization factor governing this effect and is capable of changing its sign. The second is an effective interfacial exchange coupling $-J\cos\theta$ between the ferromagnets. One finds that the magnitude and *sign* of J depend on the height of the barrier and the Stoner splitting in the ferromagnets. The third is a new, irreversible exchange term in the coupled dynamics of the ferromagnets. For one sign of external voltage V, this term describes relaxation of the Landau-Lifshitz type. For the opposite sign of V, it describes a pumping action which can cause spontaneous growth of magnetic oscillations. All of these effects were investigated consistently by analyzing the transmission of charge and spin currents flowing through a rectangular barrier separating free-electron metals. In application to Fe-C-Fe junctions, the theory predicts that the valve effect is weak and that the coupling is antiferromagnetic ($J<0$). Relations connecting the three effects suggest experiments involving small spatial dimensions.

I. INTRODUCTION

This paper treats three related phenomena involving two ferromagnets separated by a nonmagnetic tunneling barrier. Each phenomenon involves the relative directions of the two magnetic vectors. First is the magnetic *valve*, in which the tunneling conductance $G=G_0(1+\epsilon\cos\theta)$ varies with angle θ between the magnetic vectors. The first report of this effect cited $\epsilon=0.07$ for a Fe-Ge-Co film junction.[1] Smaller valve effects were observed when the barrier was antiferromagnetic NiO rather than diamagnetic Ge.[2,3]

The second effect is the Heisenberg-like interfacial exchange coupling energy $-J\cos\theta$. Ferromagnetic-resonance evidence[4,5] of exchange coupling in Fe-C-Fe films supports this concept if the 10–20 Å of amorphous carbon barrier can be regarded as insulating.

The third effect, newly predicted here, is a dissipative interfacial exchange interaction proportional to external voltage V. For one sign of V, this term describes Landau-Lifshitz-like relaxation of magnetic vibration. For the opposite sign of V, it describes a pumping action which, if strong enough, can excite free vibration.

Our treatment relies on the free-electron model of the conduction electrons. Besides being simple, this model is favored by M. B. Stearns' theory (see Sec. VI) of spin-polarized tunneling between iron-group ferromagnetic metals and superconductors.[6] These studies indicate that tunneling through Al_2O_3 film barriers originates or terminates in strongly conducting bands which are partially polarized by exchange coupling to weakly conducting strongly polarized 3D bands. In rare-earth metals, whose spin density is very localized in $4f$ states, polarized tunneling[6] even more strictly involves conduction bands which ought to resemble those of free electrons. (Spin-polarized *photoemission*, on the other hand, involves d electrons with large state density which are less well represented by free electrons.[6])

We proceed systematically with analysis of (Sec. II) spin and charge transmission coefficients, (Sec. III) the magnetic valve effect, (Sec. IV) energy-conserving exchange coupling, (Sec. V) dissipative exchange coupling, and conclude with (Sec. VI) a discussion of results.

II. TRANSMISSION COEFFICIENTS

Consider two ferromagnetic conductors separated by a plane nonmagnetic tunneling barrier. The barrier may be provided by a thin ($\sim 5-100$ Å) insulating film or vacuum. In a free-electron approximation of the spin-polarized conduction electrons inside each ferromagnet, the longitudinal part of the effective one-electron Hamiltonian may be written

$$\mathcal{H}_\xi=-\tfrac{1}{2}(d/d\xi)^2+U(\xi)-\mathbf{h}(\xi)\cdot\boldsymbol{\sigma}\ . \tag{2.1}$$

Here our system of units incorporates unit electron mass and unit Planck constant. Equation (2.1) includes terms due to kinetic energy $-(\frac{1}{2})(d/d\xi)^2$, potential $U(\xi)$ and internal exchange energy $-\mathbf{h}\cdot\boldsymbol{\sigma}$ where $-\mathbf{h}(\xi)$ is the molecular field and $\boldsymbol{\sigma}$ ($=2\mathbf{s}$) is the conventional Pauli spin operator. Although transverse momentum $\mathbf{k}_\parallel$ is omitted from the above notations, the *effects* of summation over $\mathbf{k}_\parallel$ will be accounted for in our results.

In our model of tunneling, we provisionally consider a vanishing external voltage V, and use the rectangular barrier $U=U_0$ for $0<\xi<d$ and $U=0$ otherwise, as indicated in Fig. 1. By assumption, $\mathbf{h}=0$ inside the barrier. But $\mathbf{h}=\mathbf{h}_A$ or $\mathbf{h}_B$ is constant, with $|\mathbf{h}_A|=|\mathbf{h}_B|=h_0$, within each semi-infinite ferromagnet; the two ferromagnets, A and B, have identical material properties except when ex-

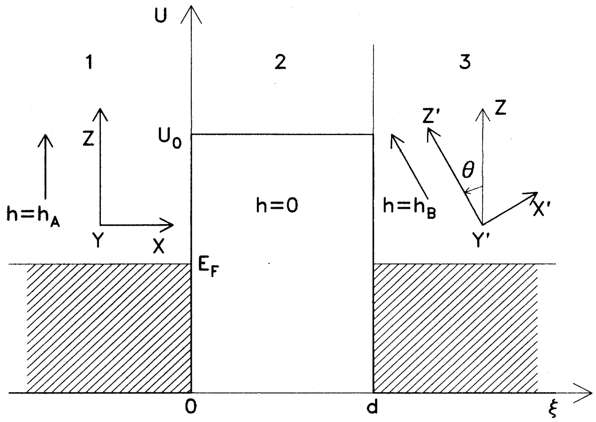

FIG. 1. Schematic potential for two metallic ferromagnets separated by an insulating barrier. The molecular fields $\mathbf{h}_A(t)$ and $\mathbf{h}_B(t)$ within the magnets form angle θ. They are instantaneously parallel to the static axes z and z' at $t=0$.

plicitly noted. However, the directions of $\mathbf{h}_A$ and $\mathbf{h}_B$, as well as the corresponding spin quantization axes z and z', differ by angle θ (see Fig. 1). Note that only the mutual relationship between coordinate systems x,y,z and x',y',z' matters. Their orientation with respect to the plane of the junction does not matter. Inside the ferromagnets, the one-electron energy is

$$E_\xi=\tfrac{1}{2}k_\sigma^2-\sigma\mathbf{h}_0, \quad \sigma=\pm 1 , \tag{2.2}$$

where k_σ is electron momentum. (Henceforth we use the notations $\sigma=\pm 1$ and $\sigma=\uparrow,\downarrow$ interchangeably.) The density of states $\rho_{\uparrow\downarrow}$ has the schematic form shown in Fig. 2. Inside the barrier, the energy is

$$E_\xi=-\tfrac{1}{2}\kappa^2+U_0, \quad \sigma=\pm 1 , \tag{2.3}$$

where $i\kappa$ is imaginary electron momentum. A similar Hamiltonian, but with $h\neq 0$ only inside the barrier, has

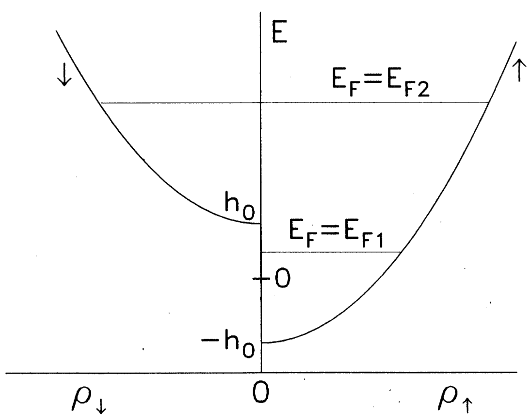

FIG. 2. Density of spin-up ($\rho_\uparrow$) and spin-down ($\rho_\downarrow$) electrons, showing position of Fermi energy E_F for one-band E_{F1} and two-band E_{F2} models of a ferromagnet.

been investigated in connection with the traversal time of tunneling.[7]

Our procedure is this: In zero-order approximation, $d=\infty$ and the magnetic states are quiescent. The total spins $\mathbf{S}_A$, $\mathbf{S}_B$ (parallel to $\mathbf{h}_A$ and $\mathbf{h}_B$, respectively) per unit junction area are not strictly defined for semi-infinite magnets. Nevertheless, we take d finite but large, and calculate the lowest-order contributions to the areal charge-current density I_e and spin-current density $\mathbf{I}_S=d\mathbf{S}_B/dt=-d\mathbf{S}_A/dt$ at the instant $t=0$. Both I_e and $\mathbf{I}_S$ will be evaluated using stationary wave functions in the same spirit as in the elementary theory of tunneling resistance. We will thus find included within $\mathbf{I}_S$ both energy conserving (Heisenberg-type) and novel nonconserving interfacial exchange terms. [We will use the fact that $\mathbf{S}_A(t)$ and $\mathbf{S}_B(t)$ can vary with t only by flow through the barrier because our assumed N-electron Hamiltonian $\mathcal{H}_N$ has only Coulomb and kinetic terms which commute with $\mathbf{S}_A$ and $\mathbf{S}_B$ when the electrodes are separated. In the absence of explicit spin-dependent terms, e.g., spin-orbit coupling or Zeeman energy, internal contributions to $\dot{\mathbf{S}}_A$ and $\dot{\mathbf{S}}_B$ must vanish.] The coordinate systems x,y,z and x',y',z' are considered *fixed* and therefore valid only in a small neighborhood of $t=0$, since $\mathbf{h}_A$ and $\mathbf{h}_B$ will generally depend on t.

Consider a spin-up incident plane wave having unit incident particle flux in region 1 (ferromagnet $A,\xi<0$ in Fig. 1). Considering the effects of all boundary conditions, the eigenfunction of $\mathcal{H}_\xi$ (eigenvalue E_ξ) in region 1 has spin components

$$\psi_{\uparrow 1}=k_\uparrow^{-1/2}e^{ik_\uparrow\xi}+R_\uparrow e^{-ik_\uparrow\xi}, \quad \psi_{\downarrow 1}=R_\downarrow e^{-ik_\downarrow\xi} , \tag{2.4}$$

where $R_\uparrow, R_\downarrow$ are coefficients to be determined. In region 2 (barrier, $0\le\xi\le d$), the wave is evanescent and has components

$$\psi_{\sigma 2}=A_\sigma e^{-\kappa\xi}+B_\sigma e^{\kappa\xi}, \quad \sigma=\uparrow,\downarrow , \tag{2.5}$$

with A_σ and B_σ to be determined. Region 3 (ferromagnet B, $d\le\xi$) has only the transmitted wave

$$\psi'_{\sigma 3}=C_\sigma e^{ik_\sigma(\xi-d)}, \quad \sigma=\uparrow,\downarrow , \tag{2.6}$$

with C_σ to be determined. The axis of spin quantization for the above waves is z (direction $\mathbf{h}=\mathbf{h}_A$) in regions 1 and 2, and z' (direction $\mathbf{h}=\mathbf{h}_B$) in region 3. Note that we need consider only real values of $k_\uparrow$ and κ. The wave vector $k_\downarrow=ip$ $(0\le p<\infty)$ for $E_\xi=E_{F1}$ (Fermi energy) is imaginary in the one-band case (see Fig. 2). It is real in the two-band case (see Fig. 2), $E_\xi=E_{F2}>h_0$, and limited to the range $0<k_\downarrow<k_\uparrow$.

To complete the solution of the Schrödinger equation, one must find the eight unknowns $R_\sigma, A_\sigma, B_\sigma, C_\sigma$ $(\sigma=\uparrow,\downarrow)$ by matching ψ_σ and $d\psi_\sigma/d\xi$ at the interfaces $\xi=0$ and $\xi=d$. The change in quantization axis at $\xi=d$ requires the spinor transformation

$$\psi_{\uparrow 2}=\psi'_{\uparrow 3}\cos(\theta/2)+\psi'_{\downarrow 3}\sin(\theta/2) , \tag{2.7a}$$

$$\psi_{\downarrow 2}=-\psi'_{\uparrow 3}\sin(\theta/2)+\psi'_{\downarrow 3}\cos(\theta/2) , \tag{2.7b}$$

and similarly for the derivatives.

Some algebra produces the following approximate solution for the coefficients, which is sufficiently accurate to calculate transmissivity to leading order in $e^{-\kappa d}$:

$$R_\sigma = A_\sigma + B_\sigma - k_\uparrow^{-1/2}\delta_{\sigma,1} \quad (\sigma=\uparrow,\downarrow) , \tag{2.8a}$$

$$A_\uparrow = \frac{2k_\uparrow^{1/2}}{(k_\uparrow + i\kappa)} + \frac{(\kappa + ik_\uparrow)B_\uparrow}{(\kappa - ik_\uparrow)} , \tag{2.8b}$$

$$B_\uparrow = \frac{-2ik_\uparrow^{1/2}[\kappa^2 + k_\uparrow k_\downarrow + i\kappa(k_\uparrow - k_\downarrow)\cos\theta]e^{-2\kappa d}}{(\kappa - ik_\uparrow)^2(\kappa - ik_\downarrow)} , \tag{2.8c}$$

$$A_\downarrow = \frac{(\kappa + ik_\downarrow)B_\downarrow}{\kappa - ik_\downarrow} , \tag{2.8d}$$

$$B_\downarrow = \frac{2\kappa k_\uparrow^{1/2}(k_\downarrow - k_\uparrow)e^{-2\kappa d}\sin\theta}{(\kappa - ik_\uparrow)^2(\kappa - ik_\downarrow)} , \tag{2.8e}$$

$$C_\uparrow = \frac{-4ik_\uparrow^{1/2}\kappa e^{-\kappa d}\cos(\theta/2)}{(\kappa - ik_\uparrow)^2} \tag{2.8f}$$

$$C_\downarrow = \frac{-4ik_\uparrow^{1/2}\kappa e^{-\kappa d}\sin(\theta/2)}{(\kappa - ik_\downarrow)(\kappa - ik_\uparrow)} . \tag{2.8g}$$

We write without proof the expressions for spin transmissivity

$$T_z = \mathrm{Im}\sum_\sigma \sigma\psi_\sigma^* \left[\frac{d\psi_\sigma}{d\xi}\right] , \tag{2.9}$$

$$T_+ \equiv T_x + iT_y = i\left[\frac{d\psi_\uparrow^*}{d\xi}\psi_\downarrow - \psi_\uparrow^*\frac{d\psi_\downarrow}{d\xi}\right] . \tag{2.10}$$

where $\mathbf{T}=(T_x,T_y,T_z)$ is the expectation value of Pauli spin ($\boldsymbol{\sigma}=2\mathbf{s}$) transmitted through the plane with given ξ. Note that the above expression for T_z becomes the conventional particle transmissivity when the factor σ in the summand is removed:

$$T_p = \mathrm{Im}\sum_\sigma \psi_\sigma^* \frac{d\psi_\sigma}{d\xi} . \tag{2.11}$$

Summations of $-eT_p$ and $\hbar\mathbf{T}/2$ over occupied states give the total charge (I_e) and spin ($\mathbf{I}_S$) currents per unit area flowing from A to B. By continuity, I_e and $\mathbf{I}_S$ do not depend on ξ. However, the molecular field $\mathbf{h}\neq 0$ within the ferromagnets causes irrelevant internal spin changes ($d\mathbf{T}/d\xi\neq 0$), whose sum over E_ξ and $\mathbf{k}_\parallel$ necessarily vanishes. Even states of small E_ξ having negligible barrier penetration would have to be included. Therefore, in these sums, $\mathbf{T}$ and T_e are conveniently evaluated anywhere within the barrier ($0\le\xi\le d$), where $\mathbf{h}=0$ and $d\mathbf{T}/d\xi=0$. Only states with E_ξ near E_F contribute appreciably when d is large.

By substitution of Eqs. (2.8b)–(2.8e) into (2.5) and (2.9)–(2.11) at $\xi=0^+$, one readily finds to first order in $e^{-2\kappa d}$,

$$T_p = T_z = \mathrm{Im}(a + ib\cos\theta), \quad T_+ = -b\sin\theta , \tag{2.12}$$

with $a = AC$, $b = BC$, and

$$A = \kappa^2 + k_\uparrow k_\downarrow, \quad B = \kappa(k_\uparrow - k_\downarrow) , \tag{2.13}$$

$$C = \frac{8\kappa k_\uparrow e^{-2\kappa d}}{(\kappa^2 + k_\uparrow^2)(\kappa - ik_\uparrow)(\kappa - ik_\downarrow)} . \tag{2.14}$$

These relations are valid for ↑ incident from the left. The other cases, ↑ incident from the right, and ↓ incident from left or right, are obtainable from them by symmetry transformations. The approximate relation $T_p = T_z$ follows from the relation $|\psi_{\downarrow 2}|/|\psi_{\uparrow 2}| \sim e^{-2\kappa d}$ which allows neglect of the $\sigma=-1$ contributions in Eqs. (2.9) and (2.11).

III. MAGNETIC VALVE

We will consider only the absolute zero of temperature. For a small external voltage V, one may use the $V=0$ wave functions of the previous section. In the limit of a small barrier factor $e^{-2\kappa d}$, a narrow distribution of electrons near normal incidence and with E_ξ near E_F carry most of the current. Therefore we may replace $\kappa(E_\xi)$ and $k_\sigma(E_\xi)$ with $\kappa(E_F)$ and $k_\sigma(E_F)$, respectively, in calculating the surface conductance G due to tunneling. This replacement will be understood in all expressions that follow except where noted otherwise. By summing the charge transmission over E_ξ and $\mathbf{k}_\parallel$ for occupied states in the usual manner,[8] one finds the conventional expression, in our notation,

$$I_e/V = G = (e^2/8\pi^2\hbar)(\kappa T_p/d) \tag{3.1}$$

for areal current density I_e, considering one initial spin direction, where e is the electron charge. We have two cases, one-band and two-band (see Fig. 2). For the one-band case $k_\downarrow = ip$ is imaginary. Writing $a = a' + ia''$, $b = b' + ib''$, we find

$$a'' = b' = \frac{8\kappa^2 k_\uparrow^2 e^{-2\kappa d}}{(\kappa^2 + k_\uparrow^2)^2} , \tag{3.2}$$

and the first of Eqs. (2.12) becomes

$$T_p = 8\kappa^2 k_\uparrow^2 e^{-2\kappa d}(1+\cos\theta)/(\kappa^2 + k_\uparrow^2)^2, \quad (k_\downarrow = ip) . \tag{3.3}$$

This substituted into (3.1) gives G for the one-band case. It represents a *perfect* magnetic valve in the sense that $G=0$ for $\theta=\pi$. Note that G does not depend on p, and therefore not on h_0 for given $U_0 - E_F$.

In the two-band case, $k_\downarrow$ is real. Now electrons with both values of σ are incident and T_p in Eq. (3.1) must be replaced by $T_{p\uparrow} + T_{p\downarrow}$. Here $T_{p\uparrow}$ is given by Eqs. (2.12)–(2.14) and $T_{p\downarrow}$ is given by the same expression with $k_\uparrow$ and $k_\downarrow$ interchanged. (Note that now $a''\neq b'$.) The result is now an *imperfect* magnetic valve with conductance of the form

$$G = G_{fbf}(1 + P_{fb}^2\cos\theta), \quad |P_{fb}| \le 1 , \tag{3.4}$$

where the effective *spin polarization* of the ferromagnetic-barrier couple is

$$P_{fb} = \frac{(k_\uparrow - k_\downarrow)}{(k_\uparrow + k_\downarrow)}\frac{(\kappa^2 - k_\uparrow k_\downarrow)}{(\kappa^2 + k_\uparrow k_\downarrow)} . \tag{3.5}$$

The mean surface conductance is now

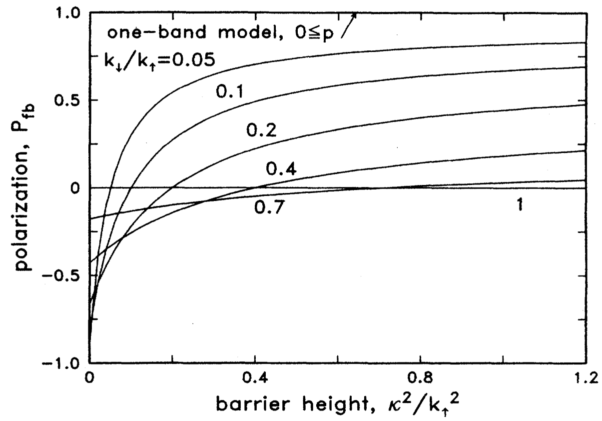

FIG. 3. Effective spin polarization P_{fb} for the magnetic valve effect, Eqs. (3.4) and (3.5). For Fe, Co, Ni and their alloys, the barrier height in eV is about 5 times the abcissa $\kappa^2/k_\uparrow^2$.

$$G_{fbf}=\frac{\kappa}{\hbar d}\left[\frac{e\kappa(\kappa^2+k_\uparrow k_\downarrow)(k_\uparrow+k_\downarrow)}{\pi(\kappa^2+k_\uparrow^2)(\kappa^2+k_\downarrow^2)}\right]^2 e^{-2\kappa d} \quad (3.6)$$

for the two-band case.

In a more general treatment of this problem, the ferromagnetic electrodes f and f' have different compositions. Thus the quantities $k_\uparrow$ and $k_\downarrow$ are assumed to be different for the two electrodes. One finds easily,

$$G=G'_{fbf'}(1+P_{fb}P_{f'b}\cos\theta)\ , \quad (3.7)$$

where P_{fb} ($P_{f'b}$) is given by Eq. (3.5) with $k_\uparrow,k_\downarrow$ replaced by $k_{\uparrow f},k_{\downarrow f}$ ($k_{\uparrow f'},k_{\downarrow f'}$), respectively. Defining $P_{fb}=1$ for a one-band magnet determines the sign convention for the two-band case as given in Eq. (3.5). A plot of P_{fb} versus reduced barrier height for various $k_\downarrow/k_\uparrow$ appears in Fig. 3.

IV. CONSERVATIVE EXCHANGE COUPLING

We now analyze the spin current to obtain exchange coupling effects. First we consider the case $V=0$, in which there is no electricity flowing through the barrier and no energy dissipation. The symmetry of identical ferromagnets now implies some simplification of spin-transmission effects. Let the spinor rate $d\sigma_A/dt$ combine $A\rightarrow B$ transmission originating from one incident ↑ (z-axis quantization) wave in magnet A ($-\mathbf{T}$ evaluated above) with the corresponding $B\rightarrow A$ transmission originating from one ↑ (z'-axis quantization) wave with the same E_ξ in magnet B. By means of the appropriate transformation connecting frames x,y,z and x',y',z' one readily finds from Eq. (2.12) the relations (in frame x,y,z)

$$d\sigma_A/dt=[(b'-a'')\sin\theta, 2b''\sin\theta\ ,(b'-a'')(1-\cos\theta)]\ . \quad (4.1)$$

For the one-band case ($k_\downarrow=ip$), the equality (3.2) causes this to reduce to

$$d\sigma_A/dt=(0,2b''\sin\theta,0)\ . \quad (4.2)$$

(Of course, we knew that $\dot\sigma_{Az}$ must vanish because of the equality $T_z=T_p$ and the fact that the particle transmissions much cancel by symmetry in case $V=0$.)

The surviving component $\dot\sigma_{Ay}$, whose direction and angular dependence are given by $\mathbf{h}_A\times\mathbf{h}_B$, represents effective spin precession associated with an effective Heisenberg-like interface coupling. It must be summed over occupied states to obtain the spin-current component $-I_y=\dot S_{Ay}$. We can infer the result from the corresponding electric-current expression (3.1). In the one-band case we simply replace I_e with $I_{Sy}\equiv-\dot S_{Ay}$, $-eT_e$ with $\hbar\dot\sigma_{Ay}/2$ evaluated at $E_\xi=E_F$, and eV with

$$W=\int_{E=E_{\min}}^{E=E_F}\exp\{2d[\kappa(E_F)-\kappa(E_\xi)]\}dE_\xi$$
$$\cong\kappa/2d=(U_0-E_F)/\kappa d \quad (4.3)$$

evaluated in the limit of large κd. Using (4.2) we thus find

$$J=\dot S_{Ay}/\sin\theta=(U_0-E_F)b''/8\pi^2d^2\ , \quad (4.4)$$

where we regard $\dot S_{Ay}$ as an effective precession of moment A due to an effective areal Heisenberg coupling energy $-J\cos\theta$. Substituting the one-band condition $k_\downarrow=ip$, into Eqs. (2.13)–(2.14), we find

$$b''=\frac{8\kappa^2k_\uparrow(k_\uparrow^2-\kappa p)}{(\kappa^2+k_\uparrow^2)^2(\kappa+p)}e^{-2\kappa d}\ , \quad (4.5)$$

to complete the formula (4.4) for J. The J-proportional dimensionless factor $(\kappa^2/k_\uparrow^2)b''e^{2\kappa d}$ is plotted versus reduced barrier height $\kappa^2/k_\uparrow^2$ in Fig. 4.

In the two-band case we must replace our wave pair with a quartet to include effects of incident ↓ waves. In Eqs. (2.12)–(2.14) one needs to interchange $k_\uparrow$ and $k_\downarrow$ to find the corresponding transmission coefficients, being mindful of the rotation in coordinate axes. For the quartet one finds the relation (4.2) holding with $k_\uparrow$ and $k_\downarrow$

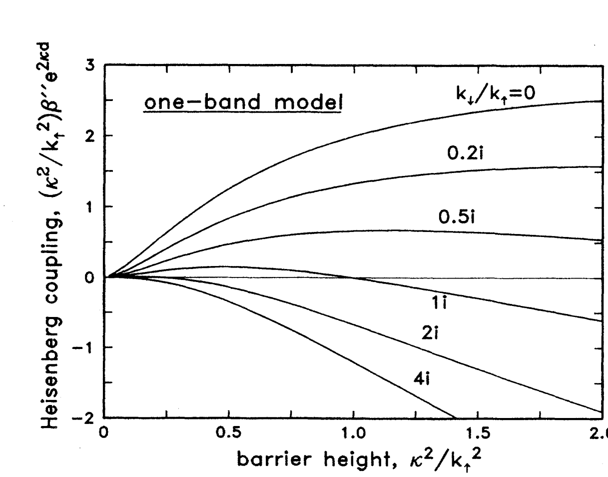

FIG. 4. Reduced Heisenberg-exchange coupling $(\kappa^2/k_\uparrow^2)\beta''e^{2\kappa d}$ vs reduced barrier height in the one-band model, Eq. (4.5).

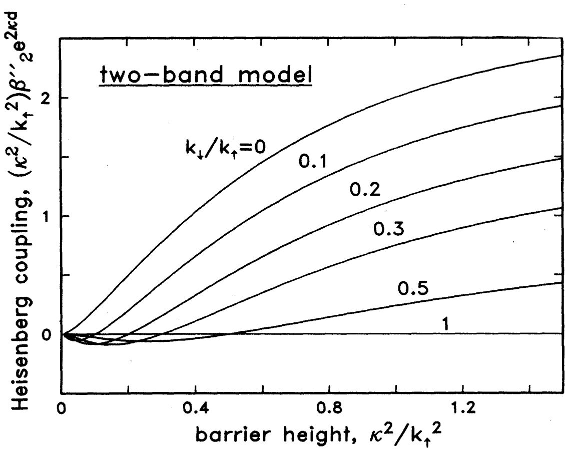

FIG. 5. Reduced Heisenberg-exchange coupling $(\kappa^2/k_\uparrow^2)\beta_2'' e^{2\kappa d}$ vs barrier height in the two-band model, Eq. (4.6).

real and b'' replaced by

$$b_2''=\frac{8\kappa^3(\kappa^2-k_\uparrow k_\downarrow)(k_\uparrow-k_\downarrow)^2(k_\uparrow+k_\downarrow)}{(\kappa^2+k_\uparrow^2)^2(\kappa^2+k_\downarrow^2)^2}e^{-2\kappa d} \quad (k_\uparrow>k_\downarrow>0)\ . \tag{4.6}$$

With this replacement, Eq. (4.4) gives the coupling for the two-band case. The proportional factor $(\kappa^2/k_\uparrow^2)b_2''e^{2\kappa d}$ is plotted in Fig. 5.

One impediment to comparing tunneling theory with experiment is the great sensitivity of the barrier factor $e^{-2\kappa d}$. It therefore may be useful to measure the magnetic valve and exchange effects in the same specimen. One can then relate the two effects, thus eliminating this sensitive factor. If we take the liberty of applying our semi-infinite-magnet results to a *finite* system in which the magnet A has a finite total number of N_A of uncompensated spins, then we may define the effective exchange field $H_{AB}=J/m_A$ of B acting on A where m_A is the moment of magnet A per unit junction area. Applying the one-band equations (3.6), (4.4), and (4.5), we find the relation

$$H_{AB1}=\frac{\hbar(U_0-E_F)}{\beta R_0 N_A e^2}\frac{(k_\uparrow^2-\kappa p)}{\kappa k_\uparrow(\kappa+p)d}\ , \tag{4.7}$$

where β is the Bohr magneton and R_0^{-1} is G_0 times the junction area. The factor $e^{-2\kappa d}$ is absent from this relation and only algebraic dependences on κ and d remain. For the two-band case the corresponding relation is

$$H_{AB2}=\frac{\hbar(U_0-E_F)}{\beta R_0 N_A e^2}\frac{(\kappa^2-k_\uparrow k_\downarrow)(k_\uparrow-k_\downarrow)^2}{d(\kappa^2+k_\uparrow k_\downarrow)^2(k_\uparrow+k_\downarrow)}\ . \tag{4.8}$$

V. DISSIPATIVE EXCHANGE COUPLING

In the presence of voltage $V\neq 0$, the tunneling contributions to $\dot S_{Ax}$ and $\dot S_{Az}$ will no longer vanish. According to Eq. (2.12), the relevant single-wave spin-transmission coefficients are

$$T_x=-b'\sin\theta,\quad T_z=a''+b'\cos\theta\ . \tag{5.1}$$

The contribution of tunneling to $\dot S_{Az}$, which we term *longitudinal*, may be neglected because, if uncompensated by other effects, it would represent change of the magnetic-order parameter. However, in reality, spin-orbit and dipolar processes internal to the magnet cause the order parameter to relax very rapidly toward equilibrium. Moreover, we envision a closed electrical circuit in which spin-polarized electrons are exchanged through electrical contacts with paramagnetic conductors. Indeed, in the simple one-band case the latter mechanism alone would guarantee the equation $\dot S_{Az}=0$ because it is equivalent to charge neutrality in a one-band ferromagnet. Effects of this so-called *spin injection* between ferro- and paramagnets have been reported recently.[9]

Thus the tunneling contribution to the component $\dot S_{Ax}$, which we term *transverse*, describes the principal dynamical consequence of dissipative exchange. In the case $V\equiv V_B-V_A>0$, electrons flow from A to B. In the one-band case we substitute I_{Sx} for I_e and $\frac{1}{2}\hbar T_x$ for $-eT_p$ in Eq. (3.1), to find

$$-I_{Sx}\equiv\dot S_{Ax}=(e\kappa b'/16\pi^2 d)(V_B-V_A)\sin\theta\ , \tag{5.2}$$

where b' is given by Eq. (3.2).

To complete the discussion, we ought to consider also the opposite case $V_A>V_B$ in which the electron flow is reversed. It is easier, instead, to look at the tunneling contribution to $\dot{\mathbf S}_B$ keeping $V_B>V_A$. According to our argument above, the longitudinal component $\dot S_{Bz'}$ is less interesting. For the transverse effect we need only

$$T_{x'}=T_x\cos\theta+T_z\sin\theta=a''\sin\theta\ , \tag{5.3}$$

with the last equality due to Eq. (5.1). Since we have $a''=b'$ in the one-band case, $\dot S_{Ax}=-\dot S_{Bx'}$ follows from Eqs. (5.1) and (5.3).

For the one-band case, we may rewrite the transverse exchange effects in the coordinate-free form

$$\dot{\mathbf S}_{At}=D(V_B-V_A)\hat{\mathbf S}_A\times(\hat{\mathbf S}_A\times\hat{\mathbf S}_B)\ , \tag{5.4}$$

where, $\wedge$ designates unit vector, and from Eqs. (3.2) and (5.2),

$$D=\frac{e\kappa^3 k_\uparrow^2 e^{-2\kappa d}}{2\pi^2 d(\kappa^2+k_\uparrow^2)^2}\ . \tag{5.5}$$

Interchanging the subscripts $A\leftrightarrow B$ in Eq. (5.4) gives $\dot{\mathbf S}_{Bt}$.

Now consider the two-band case in which one must combine $k_\uparrow$ and $k_\downarrow$ incident waves. To calculate $\dot S_{Ax}$ for $V_B>V_A$, we must write $T_x-T_x(k_\uparrow\leftrightarrow k_\downarrow)$ in place of T_x, with the minus sign due to 180° rotation of xyz axes about y. We find Eq. (5.4) again but now with D replaced by

$$D_2=\frac{e\kappa^3(\kappa^4-k_\uparrow^2k_\downarrow^2)(k_\uparrow^2-k_\downarrow^2)}{2\pi^2 d(\kappa^2+k_\uparrow^2)^2(\kappa^2+k_\downarrow^2)^2}e^{-2\kappa d}\ . \tag{5.6}$$

For $V_B>V_A$, $\dot{\mathbf S}_{At}$ and $\dot{\mathbf S}_{Bt}$ have the directions indicated in the planar Fig. 6. A heuristic explanation of these precession directions follows. Since $\hat{\mathbf S}_A$-polarized electrons impinge on magnet B, surely $\hat{\mathbf S}_B$ must relax *toward* $\hat{\mathbf S}_A$

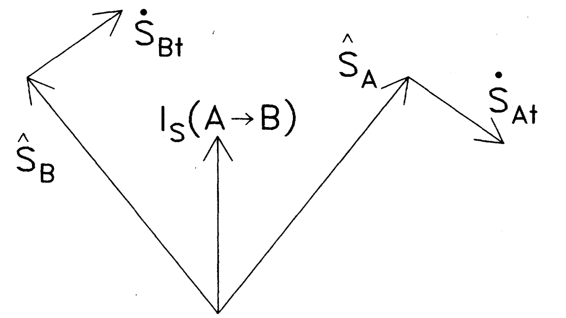

FIG. 6. Scheme of spin-vector dynamics due to the transverse terms of dissipative exchange coupling induced by an external voltage across the barrier.

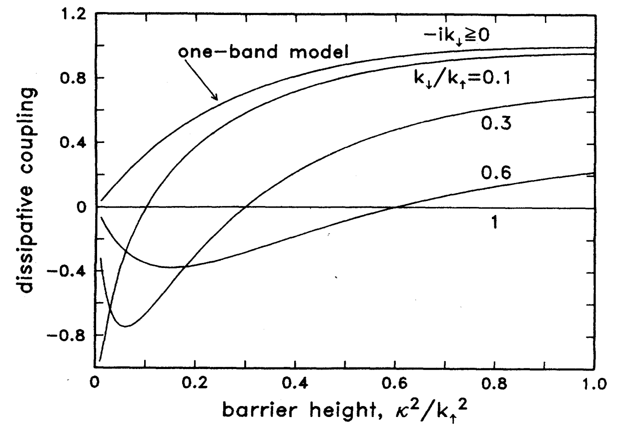

FIG. 7. Dissipative exchange coupling $De^{2\kappa d}$ (in arbitrary units) vs barrier height, Eq. (5.6).

(Fig. 6). Using optical parlance, B is the "analyzer" which prefers to receive electrons polarized in direction $\hat{\mathbf{S}}_B$. Thus magnet A is *losing* more-or-less B-polarized electrons and $\hat{\mathbf{S}}_A$ must drift *away* from $\hat{\mathbf{S}}_B$, as shown.

One can show that actually the spin current through the junction (including both longitudinal and transverse terms) is given simply by

$$\mathbf{I}_S(A\to B)=D(V_B-V_A)(\hat{\mathbf{S}}_A+\hat{\mathbf{S}}_B)\ , \tag{5.7}$$

with D given by Eq. (5.5) and (5.6). Its direction takes the mean of $\hat{\mathbf{S}}_A$ and $\hat{\mathbf{S}}_B$ (see Fig. 6). One thus has the trivial relation

$$\dot{\mathbf{S}}_{At}=-\mathbf{I}_S+(\mathbf{I}_S\cdot\hat{\mathbf{S}}_A)\hat{\mathbf{S}}_A \tag{5.8}$$

satisfied by the expressions (5.4) and (5.7).

VI. DISCUSSION

The preceding elementary model does not take account of generally important complications such as barrier nonuniformity, crystal potential, electron-electron correlations, spin waves, and inelastic tunneling. However, it does provide a basis for initial appraisal of exchange effects arising from tunneling between ferromagnets.

Notable is the sensitivity of our results, for given κd, to the height of the tunneling barrier ($\propto\kappa^2$) above E_F, as plotted in Figs. 3, 4, 5, and 7. In the one-band case, the Heisenberg-exchange expression (4.5) has the zero $\kappa=k_\uparrow^2/p$. In the two-band case the valve (3.5) and both exchange expressions (4.6) and (5.6) have the zero $\kappa^2=k_\uparrow k_\downarrow$. Thus varying κ can effect sign changes in all three effects under discussion. These results contradict the plausible notion that a quantity like P_{fb} (polarization) is characteristic of the electron structure of the *electrode alone* and would have the same value (sign at least) in any tunneling measurement. In addition, the changes of sign with variation of band splitting (h_0 or $k_\downarrow$ and p) made evident by the same zeros, is striking. Indeed, the coupling, J, exhibited in Eqs. (4.5) and (4.6) and Figs. 4 and 5, has two sign changes as h_0 varies from 0 to ∞ and both zeros are crossed. [Note the relations $k_\downarrow^2=k_\uparrow^2-4h_0$ and $p^2=4h_0-k_\uparrow^2$ from Eq. (2.2).] Thus the present simple model shows that the character of the electrode-barrier interface greatly influences exchange effects. It is likely that more detailed treatment including crystal potential, surface states and the like, would also reveal sensitivity to the interface.

The polarization coefficient $P_{fb}=P_t A_{fb}$, given by Eq. (3.5) for the two-band model, measures the contribution of magnet f and its interface with barrier b to the magnetic valve effect described by

$$G=G_{fbf'}(1+C_{fbf'}\cos\theta)\ , \tag{6.1}$$

with valve coefficient $C_{fbf'}=P_{fb}P_{f'b}$. The first factor

$$P_f=\frac{k_\uparrow-k_\downarrow}{k_\uparrow+k_\downarrow}=\frac{\rho_\uparrow-\rho_\downarrow}{\rho_\uparrow+\rho_\downarrow} \tag{6.2}$$

in P_{fb} is simply the fractional spin polarization of the state densities $\rho_\uparrow$ and $\rho_\downarrow$ at the Fermi energy of electrode f in the free-electron approximation.

The interfacial second factor

$$A_{fb}=(\kappa^2-k_\uparrow k_\downarrow)/(\kappa^2+k_\uparrow k_\downarrow)\ , \tag{6.3}$$

newly introduced here, has the range $-1<A_{fb}<1$. It arises because the penetration of ψ_σ from an electrode into the barrier depends on k_σ. Note that a sufficiently low barrier results in $A_{fb}<0$, and then the interface reverses the sign of polarization of Fermi-surface electrons penetrating from magnet f into barrier b. The possibility of reconciling within this theory existing valve measurements[1–3] with those of spin-polarized ferromagnet-to-superconductor tunneling[10] and its interpretation[11] is considered separately.[12]

The close relationships between the three effects considered here might be exploited by observing more than one of them in a single specimen. Consider the relation (4.8) between resistance and Heisenberg exchange. The band-theory analysis of polarization in magnet-to-superconductor tunneling by Stearns[11] permits evaluation of this relation for Fe-C-Fe junctions recently investigated by ferromagnetic resonance.[4,5] Stearns argues that the (Fe,Co,Ni)-Al_2O_3-Al results[10] are accounted for by two (↑ and ↓) parabolic bands having practically free-electron effective masses. For Fe and Ni, these bands represent simple approximate equivalents of certain low-mass subbands found in first-principles band calculations and in Fermi surfaces independently determined by deHaas-vanAlphen studies. The respective Fermi vectors in Fe are $k_\uparrow \simeq 1.09$ Å^{-1}, $k_\downarrow = 0.42$ Å^{-1}. (According to Stearns, $k_\uparrow$ varies little with atomic number; thus we can roughly approximate the energy barrier $U_0 - E_F$ for an insulator abutting Fe, Co, Ni, or their alloys by $\simeq 5$ eV times the abscissa $\kappa^2/k_\uparrow^2$ in Figs. 3, 4, 5, and 7.) We assume $d = 8$ Å and $2\kappa \simeq 1.0$ Å^{-1}, the latter indicated by the dependence of ferromagnetic-resonance spectra on C thickness,[4,5] implying $U_0 - E_F \simeq 1.0$ eV in C. Then Eq. (4.8) reduces to

$$H_{AB}(\text{Oe})R_0(\Omega)N_A \cong -1\times 10^{10}\ \text{Oe}\,\Omega\ , \tag{6.4}$$

where R_0 is equal to the total mean resistance of the junction an Ohms. One might consider $H_{AB} \cong -100$ Oe and $R_0 \simeq 1$ m Ω to be convenient magnitudes experimentally. One can adjust d to achieve this value of H_{AB}, principally through the sensitive factor $e^{-2\kappa d}$ in Eq. (4.6). Then, Eq. (6.4) would require magnetic film A to have dimensions of the order $N_A = 10^5 \times 10^5 \times 10$ in units of interatomic distance in order to attain $R_0 \cong 1$ mΩ. For Fe films of thickness about 20 Å, this means that a junction area of about (20 μm)^2 would be needed for the resistance measurements.

An additional prediction for Fe-C-Fe junctions is the valve-effect polarization $P_{\text{FeC}} = P_{\text{Fe}} A_{\text{FeC}} = -0.13$ from relations (6.1)–(6.3). One thus predicts a remarkably small valve coefficient $C_{\text{FeCFe}} \simeq 0.017$. Moreover, the tunneling-exchange coupling predicted by Eqs. (4.4) and (4.6) is negative.

The transverse relaxation term given by Eq. (5.4) resembles Landau-Lifshitz damping. To illustrate its possible implications, we assume that it is correct even when the thickness of magnet A (and therefore $\mathbf{S}_A$) is finite rather than infinite as in the derivation. We further assume that $\hat{\mathbf{S}}_B$ is static (for example, because of large anisotropy energy), so that the effect of spin-polarized tunneling on $\mathbf{S}_A$ is a direct relaxation toward (or away) from $\pm\hat{\mathbf{S}}_B$ (the sign accords with that of J) given by Eq. (5.2). Moreover we assume that tunneling-exchange coupling dominates the dynamics of $\mathbf{S}_A$ so that Eq. (4.4) describes precession of $\mathbf{S}_A$ about $\pm\hat{\mathbf{S}}_B$ at small-amplitude frequency $\omega_0 = |\dot{\mathbf{S}}_{Ay}/S_A \sin\theta|$. It follows that the voltage-induced relative linewidth for small-amplitude precession is $(\Delta\omega)_V/\omega_0 = \dot{S}_{Ax}/\dot{S}_{Ay}$. Considering Eqs. (3.2), (4.4), (4.5), and (5.2), we find

$$(\Delta\omega)_V/\omega_0 = \frac{(V_B - V_A)e\kappa d(\kappa+p)k_\uparrow}{2(U_0 - E_F)(k_\uparrow^2 - \kappa p)} \tag{6.5}$$

for the one-band case. For the two-band case, Eqs. (4.4), (4.6), (5.4), and (5.6) give

$$|(\Delta\omega)_V/\omega_0| = \frac{|V|e(\kappa^2 + k_\uparrow k_\downarrow)d}{2(U_0 - E_F)|k_\uparrow - k_\downarrow|}\ . \tag{6.6}$$

These expressions have the following significance. To observe voltage-pumped oscillation, one may adjust the barrier thickness so that the natural frequency ω_0 attains a value in the desired range for ferromagnetic resonance (FMR) or Brillouin scattering. Then Eq. (6.5) or (6.6) tells how much change in linewidth is induced by the voltage. (Note that the exponential barrier factor has cancelled.) If the intrinsic (caused by effects other than tunneling) linewidth $(\Delta\omega)_0$ is homogeneous (e.g., spin-lattice relaxation), then satisfying the condition $(\Delta\omega)_0 + (\Delta\omega)_V < 0$ will be sufficient to provide exponentially growing oscillations in magnet A. For Fe-C-Fe junctions $|(\Delta\omega)_V/\omega_0| \cong 4\ V$(volt) and the necessary voltage will be some millivolts for narrow resonance lines.

Excessive temperature rise due to power dissipation $\Phi = V^2 G$ (per unit junction area) could limit the possibility of observing voltage-pumped precession. A useful expression for Φ follows from the relations $|J| = |\dot{S}_{Ay}/\sin\theta| = \omega_0 S_A$ due to Eq. (4.4), and

$$|\dot{S}_{Ax}| = |VD\sin\theta| = (\Delta\omega)_V|S_{Ax}| = |(\Delta\omega)_V S_A \sin\theta|$$

due to Eq. (5.4). They allow Φ to be written the form

$$\Phi = [(\Delta\omega)_V^2 S_A/\omega_0]|GJ|/D^2 \tag{6.7}$$

in which dependence of the factor GJ/D^2 on the expression $e^{-2\kappa d}$ will cancel.

Consider the one-band case. For $\theta = \pi$, $G = 0$ from Eqs. (3.1) and (3.3) and our formula gives $\Phi = 0$, meaning that more physics would be needed to estimate the power for this case. For $\theta = 0$, we substitute Eqs. (3.1), (3.3), (4.4), (4.5), and (5.5) into (6.7) to find

$$\Phi = \frac{(\Delta\omega)_V^2 n_A}{\omega_0}\,\frac{4(U_0 - E_F)(k_\uparrow^2 - \kappa p)}{d\kappa k_\uparrow(\kappa+p)}\ , \tag{6.8}$$

where n_A is the number of ferromagnetic spins per unit junction area and where $k_\uparrow^2 > \kappa p$ to ensure equilibrium at $\theta = 0$.

For the two-band case, Eqs. (3.4), (3.6), (4.4), (4.6), and (5.6) similarly give

$$\Phi = \frac{(\Delta\omega)_V^2 n_A}{\omega_0}\,\frac{2(U_0 - E_F)(k_\uparrow + k_\downarrow)(1 \pm P_{fb}^2)}{d|\kappa^2 - k_\uparrow k_\downarrow|}\ , \tag{6.9}$$

where $\pm P_{fb}^2$ corresponds to $\theta = 0, \pi$ and P_{fb} is given in Eq. (3.5). Assuming ten atomic layers of Fe ($n_A \simeq 5\times 10^8$ μm^{-2}) and a carbon barrier, our parameters reduce Eq. (6.9) to $\Phi \simeq 10^{-10}(\Delta\omega_V)^2/\omega_0$ W/μm^2. Assuming $\omega_0 = 10^{11}$ Hz, a power of $\Phi \simeq 1$ mW/μm^2 is needed for a linewidth change amounting to $(\Delta\omega)_V = -10^{-2}\omega_0$. Brillouin scattering with a small focal spot and intermittent application of voltage might provide a means to observe voltage-excited magnetic vibration without creating too great a temperature rise.

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[1]M. Julliere, Phys. Lett. **54A**, 225 (1975).
[2]U. Gäfvert and S. Maekawa, IEEE Trans. Magn. **MAG-18**, 707 (1982).
[3]Y. Suezawa and Y. Gondō, in *Proceedings of the International Symposium on Physics of Magnetic Materials,* edited by M. Takahashi, S. Maekawa, Y. Gondō, and H. Nose (World Scientific, Singapore, 1987), p. 303.
[4]M. Pomerantz, J. C. Slonczewski, and E. Spiller, J. Appl. Phys. **61**, 3747 (1987).
[5]M. Pomerantz, J. C. Slonczewski, and E. Spiller, in *Proceedings of the International Symposium on Physics of Magnetic Materials,* edited by M. Takahashi, S. Maekawa, Y. Gondō, and N. Nose (World Scientific, Singapore, 1987), p. 64.
[6]E. L. Wolf, *Principles of Electronic Tunneling Spectroscopy* (Oxford University Press, New York 1985), Sec. 8.2.4.
[7]M. Büttiker, Phys. Rev. B **27**, 6178 (1983).
[8]C. B. Duke, *Tunneling in Solids* (Academic, New York, 1969).
[9]M. Johnson and R. H. Silsbee, Phys. Rev. B **37**, 5326 (1988).
[10]P. M. Tedrow and R. Meservey, Phys. Rev. B **7**, 318 (1973); R. Meservey, D. Paraskevopoulos, and P. M. Tedrow, J. Appl. Phys. **49**, 1405 (1978).
[11]M. B. Stearns, J. Magn. Magn. Mater. **5**, 167 (1977).
[12]J. C. Slonczewski (unpublished).