Surface relaxation of the magnetization of conduction electrons in thin metallic films

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We consider the effect exerted on the width and shape of the line of paramagnetic resonance with conduction electrons by spin-orbit scattering of the electron spins by potential centers located on one film surface. The film thickness is assumed to be less than the depth of the skin layer. The method of temperature Green's functions is used to show that the line shape is asymmetrical and depends on the orientation of the film relative to the static and radio-frequency fields. In addition, a new surface contribution, which depends on the electron mean free path and on the orientation of the static magnetic field relative to the film surface, appears in the total width. Boundary conditions that take into account these effects in the calculation of the line shape by the method of Bloch's equations are proposed for the magnetization.

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1. INTRODUCTION

The main result of the theory of surface relaxation of conduction-electron (CE) spins[1–4] can be summarized in the following manner:

1) By introducing a single parameter ε, the probability of the CE spin flip following a single collision with the surface, averaged over the incidence angles, it is possible to take into account phenomenologically the surface relaxation of the CE magnetization with the aid of a boundary condition of the type

$$2\frac{\varepsilon}{L}\mathbf{M}(\mathbf{r},t)=-\left(\mathbf{n}\frac{\partial}{\partial\mathbf{r}}\right)\mathbf{M}(\mathbf{r},t),\quad \mathbf{r}\in S, \tag{1}$$

where $\mathbf{M}(\mathbf{r},t)$ is the magnetization of the CE, $\mathbf{n}$ is a unit vector normal to the surface S and directed into the interior of the sample, L is a quantity with the dimension of length, and the coefficient 2 in the left-hand side of the equation takes into account the contribution made to the magnetization relaxation by both directions of the CE spin.

2) In films having a thickness smaller than the classical skin-layer depth, the signal of paramagnetic resonance on the conduction electrons (PRCE) is described by a Lorentz absorption line.[1,4]

3) The total width of the PRCE line acquires an additional term[1–4] in the form

$$1/T_S=4\varepsilon D/dL,$$

where D is the CE diffusion coefficient and d is the film thickness.

From the conservation law for the flux of the CE through the surface, we obtain $L=4D/v_F$,[1,2] i.e.,

$$1/T_S=\varepsilon v_F/d, \tag{2}$$

where v_F is the Fermi velocity of the CE.

To explain the increase of the total width of the PRCE with decreasing temperature, it was proposed that $L\sim d$ when $l_V \ll d$,[3,4] which is equivalent to

$$1/T_S=\varepsilon v_F l_V/d^2 \tag{3}$$

accurate to a numerical coefficient of the order of unity, where l_V is the volume mean free path of the CE.

The value of ε was determined mathematically and it was shown[5] that ε is anisotropic for certain scattering mechanisms, i.e., it depends on the orientation of the quantization axis relative to the normal to the surface.

To determine the dependence of the spin surface-relaxation time T_S on the mean free path and the effect of the anisotropy of ε on the width and shape of the PRCE line, the surface relaxation will be taken into account in the present paper on a microscopic scale. The Hamiltonian contains explicitly the operators that bring about both the spin-orbit and the momentum scattering of the CE by potential centers that are randomly distributed on one of the surfaces as well as inside the volume.

It is shown by the method of temperature Green's functions that the total PRCE width line includes, simultaneously with a term of the type (2), a new contribution due to the interaction of the CE spins with the surface. This contribution depends on the mean free path and on the orientation of the static magnetic field relative to the surface; on the other hand, if identical force centers are located on both sides of the film surface, this contribution vanishes.

In addition, the anisotropy of ε leads to an asymmetry of the observed PRCE signal; this asymmetry depends on the orientation of the film relative to the direction of the radio-frequency (RF) field. Boundary conditions that take the obtained effects into account are proposed for the magnetization.

2. THEORETICAL MODEL. THE HAMILTONIAN

A metallic film is regarded as a system consisting of free electrons in a jelly of positive charge with area S ($S\to\infty$), located in the region $z_B \leqslant z \leqslant d-z_B$, where z_B

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is the width of the double layer and satisfies the condition of electric neutrality of the sample far from the surface.[6] From the electron theory of surface properties of metals[7] it is known that in a number of cases the density of the CE near the surface is described adequately by a model in which the effective single-electron potential acting on the CE is chosen in the form of a rectangular step with infinite jumps at the points $z=0$ and $z=d$. Then the wave functions and the probability density of the CE flux on the surface $(\mathbf{j}\cdot\mathbf{n})$ have the well-known form[6,5]

$$\psi_{\mathbf{k}}(\mathbf{r})=(2/Sd)^{1/2}\exp(i\mathbf{p}\rho)\sin qz,\qquad (\mathbf{jn})=(v/2Sd)\cos\theta, \tag{4}$$

where $\mathbf{k}=(p_x,p_y,q)$; $p_{x,y}=(2\pi/\sqrt{S})n_{x,y}$; $q=(\pi/d)n_z$; $n_{x,y}=0,\pm1,\pm2,\ldots$; $n_z=1,2,\ldots$; $\rho=(x,y,0)$; $v=\hbar k/m^*$ is the CE velocity; $\theta=\widehat{\mathbf{kn}}$ is the angle between $\mathbf{k}$ and the normal $\mathbf{n}$ directed along the z axis, and $0\leqslant\theta\leqslant\pi/2$.

On the surface $z=0$ are located force centers with a potential

$$U_S(\rho,z)=\sum_i u_S(\rho-\rho_i,z).$$

The coordinates $\rho_i=(x_i,y_i,0)$ of the centers are randomly and independently distributed, with a density c_S. The volume also contains randomly distributed centers with a potential

$$U_V(\mathbf{r})=\sum_j u_V(\mathbf{r}-\mathbf{r}_j)$$

and with a coordinate distribution density c_V. The centers are assumed to be short-range:

$$r_Sk_F\ll1,\quad r_Vk_F\ll1,$$

where r_S and r_V are the effective radii of the surface and volume centers, respectively, and k_F is the CE momentum at the Fermi level.

The film is assumed to be thin:

$$d\ll\delta, \tag{5}$$

where δ is the classical depth of the skin layer.

In the assumed model, the Hamiltonian of the system of CE in a metallic film with impurities in the absence of an RF field is given by

$$\mathscr{H}=\mathscr{H}_0+\mathscr{H}_1+\mathscr{H}_2.$$

$\mathscr{H}_0$ is the Hamiltonian of the CE in a static magnetic field H_0:

$$\mathscr{H}_0=\sum_{\mathbf{k}\gamma}\varepsilon_{\mathbf{k}\gamma}a_{\mathbf{k}\gamma}^+a_{\mathbf{k}\gamma},\qquad \varepsilon_{\mathbf{k}\beta}=\frac{\hbar^2k^2}{2m^*}+\frac{\Omega_0}{2},\quad \varepsilon_{\mathbf{k}\alpha}=\frac{\hbar^2k^2}{2m^*}-\frac{\Omega_0}{2}, \tag{6}$$

where the spin indices $\gamma=\beta$ and $\gamma=\alpha$ denote the states with spin up and spin down, respectively, and $\Omega_0=\mu H_0/\hbar$, where μ is the Bohr magneton. It is assumed that $\Omega_c\tau\ll1$, where Ω_c and τ are the cyclotron frequency and the mean free path of the CE. The analysis of the spin surface relaxation with allowance for quantization of the orbital motion and with allowance for the magnetic surface states is a separate problem.

$\mathscr{H}_1$ is a single-particle operator describing the potential momentum scattering

$$\mathscr{H}_1=\mathscr{H}_1^V+\mathscr{H}_1^S,$$

where $\mathscr{H}_1^V$ takes into account the scattering by the volume force centers:

$$\mathscr{H}_1^V=a_V\frac{2}{Sd}\sum_{\mathbf{k}\mathbf{k}'\gamma}\sum_j\exp[-i(\mathbf{p}-\mathbf{p}')\rho_j]\sin qz_j\sin q'z_ja_{\mathbf{k}\gamma}^+a_{\mathbf{k}'\gamma}, \tag{7a}$$

and $\mathscr{H}_1^S$ describes scattering by surface impurities:

$$\mathscr{H}_1^S=a_S\frac{2}{Sd}\sum_{\mathbf{k}\mathbf{k}'\gamma}\sum_i\exp[-i(\mathbf{p}-\mathbf{p}')\rho_i]\hat{q}\hat{q}'a_{\mathbf{k}\gamma}^+a_{\mathbf{k}'\gamma}, \tag{7b}$$

$$a_V=\int u_V(\mathbf{r})\,d\mathbf{r},\quad a_S=\int_{z\geqslant0}(zk_F)^2u_S(\mathbf{r})\,d\mathbf{r},\quad \hat{q}=\frac{q}{k_F}.$$

$\mathscr{H}_2$ is a single-particle operator describing the spin-orbit impurity scattering of the CE:

$$\mathscr{H}_2=\mathscr{H}_2^V+\mathscr{H}_2^S,$$

where $\mathscr{H}_2^V$ takes into account the scattering by the volume centers:

$$\mathscr{H}_2^V=b_V\frac{2}{Sd}\sum_{\mathbf{k}\mathbf{k}'\gamma\gamma'}\sum_j\exp[-i(\mathbf{p}-\mathbf{p}')\rho_j][\cos qz_j\sin q'z_j(\hat{\mathbf{p}}',\hat{\mathbf{q}},\sigma_{\gamma\gamma'})+\sin qz_j\cos q'z_j(\hat{\mathbf{p}},\hat{\mathbf{q}}',\sigma_{\gamma\gamma'})+i\sin qz_j\sin q'z_j(\hat{\mathbf{p}},\hat{\mathbf{p}}',\sigma_{\gamma\gamma'})]a_{\mathbf{k}\gamma}^+a_{\mathbf{k}'\gamma'}, \tag{8a}$$

and the operator $\mathscr{H}_2^S$ describes the scattering by the surface centers:

$$\mathscr{H}_2^S=b_S\frac{2}{Sd}\sum_{\mathbf{k}\mathbf{k}'\gamma\gamma'}\sum_i\exp[-i(\mathbf{p}-\mathbf{p}')\rho_i]\times[\hat{q}'(\hat{\mathbf{p}}',\hat{\mathbf{q}},\sigma_{\gamma\gamma'})+\hat{q}(\hat{\mathbf{p}},\hat{\mathbf{q}}',\sigma_{\gamma\gamma'})]a_{\mathbf{k}\gamma}^+a_{\mathbf{k}'\gamma'}, \tag{8b}$$

$$b_V=\frac{\varepsilon_F}{2mc^2}\int u_V(\mathbf{r})\,d\mathbf{r},\qquad b_S=\frac{\varepsilon_F}{2mc^2}\int_{z\geqslant0}(zk_F)u_S(\mathbf{r})\,d\mathbf{r},$$

$$\hat{\mathbf{q}}=\mathbf{n}q/k_F,\quad \hat{\mathbf{p}}=\mathbf{p}/k_F,$$

where $(\hat{\mathbf{p}},\hat{\mathbf{q}},\sigma)$ is a mixed product of the vectors $\mathbf{p}$, $\mathbf{q}$, and σ; ε_F is the Fermi energy.

It must be emphasized that we have left out of $\mathscr{H}_2^S$ the term $\sim(\hat{\mathbf{p}},\hat{\mathbf{p}}',\sigma)$, which is small in the assumed model relative to the parameter (r_Sk_F). This anisotropy of the spin-orbit interaction of the CE with the surface centers is due to the smallness of the wave function in the case of glancing incidence of the CE on the surface, since the electron must fall and be reflected tangentially to the surface in order for the angular momentum $\mathbf{L}\sim\mathbf{k}\times\mathbf{k}'$ relative to the force center to be directed normal to the surface, but the wave function of the CE tends to zero in the case of glancing incidence.

3. LINE SHAPE. CALCULATION OF MAGNETIZATION

We choose in spin space an orthogonal coordinate system (X, Y, Z) with axis $Z \parallel \mathbf{H}_0$. We direct along the X axis a radio-frequency field

$$\mathbf{H}_1(t) = \mathbf{H}_1 \cos \omega t.$$

Under condition (5), the amplitude of the RF field $\mathbf{H}_1$ is constant over the film thickness. The RF field power absorbed by the sample near resonance is described by the expression[8]

$$P = \frac{\omega H_1^2 V}{4} \operatorname{Im} \chi^-(\omega), \qquad (9)$$

where $\chi^-(\omega)$ is determined from the relation

$$M^-(\omega) = V\chi^-(\omega) H_1(\omega). \qquad (10)$$

Here $H_1(\omega) = H_1/2$ and $M^-(\omega) = M_X(\omega) - iM_Y(\omega)$ are the Fourier components of the RF field and of the sample magnetization.

$M^-(\omega)$ is expressed in terms of the retarded Green's function of the spin density of the particles, which is an analytic upward continuation of the corresponding function of the discrete frequencies. The analytic continuation can be easily carried out in general form, since the scattering centers are static. Using the standard diagram technique in the field of random centers,[9] we obtain for $M^-(\omega)$ the following expression (cf.[10]):

$$M^-(\omega+i\delta) = 2\mu^2 \int \frac{dE}{2\pi i} \{[n(E+\omega) - n(E)]\Pi(E+\omega+i\delta, E-i\delta) + n(E)\Pi(E+\omega+i\delta, E+i\delta) - n(E+\omega)\Pi(E+\omega-i\delta, E-i\delta)\}\mathbf{H}_1(\omega), \qquad (11)$$

where $n(E)$ is the Fermi distribution function, and the polarization loop

$$\Pi(E, E') = \sum_{\alpha k_1 k_2 \gamma_1 \gamma_2} \mathcal{G}(\beta k, \gamma_1 k_1; E)\mathcal{G}(\gamma_2 k_2, \alpha k; E')\Gamma(\gamma_1 k_1, \gamma_2 k_2; E, E') \qquad (12)$$

is expressed in terms of the quantity Γ, which is the aggregate of the diagrams that transform a particle from one state into another in an RF field that is constant over the film thickness. The quantity Γ satisfies the integral equation[11]:

$$\Gamma(1.2; E, E') = \sigma_{1,2} + \sum_{3,4,5,6} \mathcal{U}(1.2; 3.4; E, E')\mathcal{G}(3.5; E)\mathcal{G}(6.4; E')\Gamma(5.6; E, E'), \qquad (13)$$

where $1 = \gamma_1 \mathbf{k}_1$, $2 = \gamma_2 \mathbf{k}_2, \ldots$, $\mathcal{G}(1, 2; E)$ is the CE Green's function averaged over the impurity positions, and $\mathcal{U}(1, 2; 3, 4; E, E')$ is an irreducible four-point diagram that cannot be divided into two parts by cutting only two lines corresponding to particles.

The diagram technique that describes the averages of the Green's function, the diagrams for $\mathcal{U}$ and Γ, is shown in Fig. 1. The thin solid line represents the bare Green's function $\mathcal{G}_0$, which takes the form

$$\mathcal{G}_0(\gamma_2 \mathbf{k}_2, \alpha \mathbf{k}; E-i\delta) = \delta(\gamma_2, \alpha)\delta(\mathbf{k}_2, \mathbf{k})[E - \varepsilon_{\mathbf{k}\alpha} + \varepsilon_F - i\delta]^{-1},$$

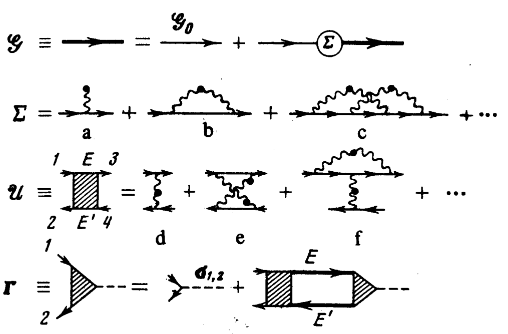

FIG. 1.

where $\delta(1, 2)$ is the Kronecker symbol, and the expression for $\varepsilon_{\mathbf{k}\alpha}$ is given in (6). The thick solid line denotes the averaged Green's function $\mathcal{G}(1, 2; E)$, which is nondiagonal with respect to the normal component of the momentum q and the spin states of the CE. This nondiagonality is due to the scattering of the CE by the surface centers and by the volume impurities located in the region of the double layer. The wavy lines correspond to matrix elements of the force-center potential, which are given in (7) and (8). The dot denotes the force center with respect to the coordinates of which averaging is carried out with weights c_V and c_S respectively for the volume and surface impurities.

The calculations will be carried out in the Born approximation, since[11] the final results, expressed in terms of the collision and relaxation times and of the parameter ε, will be valid also in general form. The simplest diagrams a and b of Fig. 1 for the self-energy part Σ correspond to the following analytic expression:

$$\Sigma(\alpha \mathbf{k}, \alpha \mathbf{k}; E-i\delta) = \Sigma_1^V + \Sigma_1^S + \Sigma_2^V + \Sigma_2^S, \qquad (14)$$

$$\Sigma_1^V = c_V a_V + \frac{i}{2\tau_V}, \quad \Sigma_1^S = \frac{c_S}{d} a_S 2\hat{q}^2 + \frac{i}{2\tau_S} 3\hat{q}^2,$$

$$\Sigma_2^V = \frac{3}{8}\frac{i}{T_V}, \quad \Sigma_2^S = \frac{\delta\Omega_0}{2} + \frac{i}{T_S}\frac{15}{8}\hat{q}^2\left(\hat{p}^2 + \frac{2}{5}\right),$$

while diagram d of Fig. 1 for the four-point function $\mathcal{U}$ corresponds to

$$\mathcal{U}(\beta \mathbf{k}, \alpha \mathbf{k}; \gamma_1 \mathbf{k}_1, \gamma_2 \mathbf{k}_1; E, E') = \mathcal{U}_1^V + \mathcal{U}_1^S + \mathcal{U}_2^V + \mathcal{U}_2^S, \qquad (15)$$

$$\mathcal{U}_1^V = [V 2\pi\rho(\varepsilon_F)\tau_V]^{-1}\delta(\beta, \gamma_1)\delta(\gamma_2, \alpha),$$

$$\mathcal{U}_1^S = [V 2\pi\rho(\varepsilon_F)\tau_S]^{-1} 9\hat{q}^2\hat{q}_1^2\delta(\beta, \gamma_1)\delta(\gamma_2, \alpha),$$

$$\mathcal{U}_2^V = [V 2\pi\rho(\varepsilon_F) 9T_V/8]^{-1}[(\hat{\mathbf{q}}, \hat{\mathbf{p}}_1, \sigma_{\beta\gamma_1})(\hat{\mathbf{q}}, \hat{\mathbf{p}}_1, \sigma_{\gamma_2\alpha}) + (\hat{\mathbf{q}}_1, \hat{\mathbf{p}}, \sigma_{\beta\gamma_1})(\hat{\mathbf{q}}_1, \hat{\mathbf{p}}, \sigma_{\gamma_2\alpha}) + (\hat{\mathbf{p}}, \hat{\mathbf{p}}_1, \sigma_{\beta\gamma_1})(\hat{\mathbf{p}}, \hat{\mathbf{p}}_1, \sigma_{\gamma_2\alpha})],$$

$$\mathcal{U}_2^S = [V 2\pi\rho(\varepsilon_F) 45T_S/4]^{-1}\hat{q}^2\hat{q}_1^2(\mathbf{n}, \hat{\mathbf{p}}+\hat{\mathbf{p}}_1, \sigma_{\beta\gamma_1})(\mathbf{n}, \hat{\mathbf{p}}+\hat{\mathbf{p}}_1, \sigma_{\gamma_2\alpha}),$$

where the indices V and S correspond to volume and surface scattering, and the indices 1 and 2 correspond to momentum and spin-orbit scattering. $V = Sd$ is the volume of the film, $\rho(\varepsilon_F) = m^* k_F / 2\pi^2\hbar^2$ is the density of states of the CE per unit volume near the Fermi level for one value of the spin. The expressions for the free-path times τ_V and τ_S, for the relaxation times T_V and T_S, and for the shift of the resonant frequency $\delta\Omega_0$ are

$$\tau_V^{-1} = c_V \frac{2\pi}{\hbar}\rho(\varepsilon_F) a_V^2, \quad \tau_S^{-1} = \frac{c_S}{d}\frac{2\pi}{\hbar}\rho(\varepsilon_F) a_S^2,$$

$$T_V^{-1} = c_V \frac{2\pi}{\hbar}\rho(\varepsilon_F) b_V^2 \frac{9}{8}, \quad T_S^{-1} = \frac{c_S}{d}\frac{2\pi}{\hbar}\rho(\varepsilon_F) b_S^2 \frac{16}{45}, \qquad (16)$$

$$\delta\Omega_0=\frac{c_S}{d}\frac{b_S}{\hbar}8\hat{q}^2(\hat{\mathbf{p}},\mathbf{n},\sigma_{\alpha\alpha}),$$

where c_V and c_S are the concentrations of the volume and surface centers, while a_V, a_S, b_V, and b_S are the constants, which enter in (7) and (8), of the interaction of the CE with the centers. We note that (16) contains in explicit form Planck's constant, which was left out before.

In the derivation of the expressions for Σ^V in (14) and $\mathcal{U}^V$ in (15) we have neglected the change of the scattering by volume centers located in the region of the double layer near the surface, since the probability of scattering by these centers is lower than the probability of scattering by all the volume centers and by the same token is lower than the surface-scattering probability.

We note that Σ_2^S has a non-zero real part that depends on the spin state, but $\operatorname{Re}\Sigma_2^V=0$. The reason is that the electric field averaged over the film (diagram a of Fig. 1) and produced by the centers situated on one surface is different from zero (is directed along the normal), so that on moving in the given electric field the CE spin energy changes by an amount $\delta\Omega_0/2$. The average electric field produced by the volume impurities is zero. It should be noted that if the identical impurities were randomly distributed on both surfaces of the film, then we would have $\operatorname{Re}\Sigma_2^S=\delta\Omega_0=0$.

The diagrams in Fig. 1, corresponding to "girths" and intersections for volume scattering, are small in comparison with diagrams without intersections in terms of the parameter $(\varepsilon_F\tau_V)^{-1}$. It was shown[12] that if individual-center scattering cross sections projected on the surface do not overlap, then the "girths" and the intersections corresponding to surface scattering can be neglected, and the off-diagonal elements of the Green's functions as well. Thus, if the conditions

$$\frac{d}{v_F}\frac{1}{\tau_S}\ll 1,\quad \frac{d}{v_F}\frac{1}{T_S}\ll 1 \tag{17}$$

are satisfied, then we can neglect in (12) and (13) the off-diagonal elements of the Green's function

$$\mathcal{G}(\gamma_2\mathbf{k}_2,\alpha\mathbf{k};E-i\delta)=\delta(\gamma_2,\alpha)\delta(\mathbf{k}_2,\mathbf{k})[E-\varepsilon_{\mathbf{k}\alpha}+\varepsilon_F-\Sigma(\alpha\mathbf{k},\alpha\mathbf{k};E-i\delta)]^{-1}, \tag{18}$$

and the ladder approximation with an irreducible four-point function $\mathcal{U}$, given graphically by diagram d of Fig. 1, is valid for the quantity $\mathbf{\Gamma}$ in (13).

It is easy to verify, by considering the perturbation-theory series for $\mathbf{\Gamma}$, that the following equality (cf.[10])

$$\int\frac{dE}{2\pi i}[n(E)\Pi(E+\omega+i\delta,E+i\delta)-n(E+\omega)\Pi(E+\omega-i\delta,E-i\delta)]=V\rho(\varepsilon_F)\sigma_{\beta\alpha}$$

is satisfied, accurate to terms small in the parameter (ω/ε_F), (Ω_0/ε_F), and $(\varepsilon_F\tau_V)^{-1}$. Taking this circumstance into account and using the usual rule for integrating a Fermi distribution function, we obtain

$$M^-(\omega+i\delta)=V\chi_0\left[\sigma_{\beta\alpha}-\omega\frac{\Pi(\omega+i\delta,0-i\delta)}{2\pi iV\rho(\varepsilon_F)}\right]\mathbf{H}_1(\omega), \tag{19}$$

where $\chi_0=2\mu^2\rho(\varepsilon_F)$ is the static susceptibility. The presence of the first term in the square brackets means that the magnetization relaxes to a local instantaneous equilibrium value.

We define the quantity $\mathbf{f}$, which is the analog of the distribution function, by the following relation:

$$\mathcal{G}(\gamma_1\mathbf{k}_1,\gamma_1\mathbf{k}_1;\omega+i\delta)\mathcal{G}(\gamma_2\mathbf{k}_1,\gamma_2\mathbf{k}_1;0-i\delta)\Gamma(\gamma_1\mathbf{k}_1,\gamma_2\mathbf{k}_1;\omega+i\delta,0-i\delta)$$
$$=[\mathcal{G}(\gamma_1\mathbf{k}_1,\gamma_1\mathbf{k}_1;\omega+i\delta)-\mathcal{G}(\gamma_2\mathbf{k}_1,\gamma_2\mathbf{k}_1;0-i\delta)]\mathbf{f}_{\gamma_1\gamma_2}(\mathbf{k}_1,\omega). \tag{20}$$

Taking (19), (12), (18), and (20) into account, we find that

$$M^-(\omega)=V\chi_0[\sigma_{\beta\alpha}-\omega\langle\mathbf{f}_{\beta\alpha}(\mathbf{k}_1,\omega)\rangle]\mathbf{H}_1(\omega), \tag{21}$$

where $\langle\mathbf{f}\rangle=\int\mathbf{f}d\Omega/2\pi$ denotes averaging over the Fermi surface. In the derivation of (21) we made use of the fact that, accurate to (ω/ε_F), (Ω_0/ε_F), and $(\varepsilon_F\tau_V)^{-1}$, the following equality is satisfied:

$$\mathcal{G}(\alpha\mathbf{k}_1,\alpha\mathbf{k}_1;0-i\delta)-\mathcal{G}(\beta\mathbf{k}_1,\beta\mathbf{k}_1;\omega+i\delta)$$
$$=2i\operatorname{Im}\mathcal{G}(\mathbf{k}_1,\mathbf{k}_1;0-i\delta)=2\pi i\delta(\varepsilon_{\mathbf{k}_1}-\varepsilon_F). \tag{22}$$

Expressing $\mathbf{f}$ from (20) in terms of $\mathbf{\Gamma}$ and using (13) we find, taking (18) and (22) into account, that f satisfies the equation

$$[\omega-\Omega_0+\Sigma(\alpha\mathbf{k},\alpha\mathbf{k};0-i\delta)-\Sigma(\beta\mathbf{k},\beta\mathbf{k};\omega+i\delta)]\mathbf{f}_{\beta\alpha}(k,\omega)$$
$$=\sigma_{\beta\alpha}+2\pi i\rho(\varepsilon_F)V\sum_{\gamma_1\gamma_2}\langle\mathcal{U}(\beta\mathbf{k},\alpha\mathbf{k};\gamma_1\mathbf{k}_1,\gamma_2\mathbf{k}_1)\mathbf{f}_{\gamma_1\gamma_2}(\mathbf{k}_1,\omega)\rangle, \tag{23}$$

where the energy dependence has been left out of the argument of the block $\mathcal{U}$. In the derivation of (23) we have neglected the elements of $\mathbf{\Gamma}$ in (13) that are not diagonal with respect to the normal component of the CE momentum, with the same accuracy as of (17). Equation (23) is a closed equation for the determination of the function f on the Fermi surface.

Owing to the symmetry of Σ relative to time reversal and owing to the fact that the average electric field produced by the exchange impurities is equal to zero, we have, as seen from (14) and (16),

$$\Sigma(\alpha\mathbf{k},\alpha\mathbf{k};0-i\delta)-\Sigma(\beta\mathbf{k},\beta\mathbf{k};\omega+i\delta)=\delta\Omega_0+2i\operatorname{Im}\Sigma(\alpha\mathbf{k},\alpha\mathbf{k};0-i\delta).$$

Since the momentum scattering is much stronger than the spin-orbit scattering, we seek a solution of (23) near resonance in the form of a series in the small parameters $(\operatorname{Im}\Sigma_2/\operatorname{Im}\Sigma_1)$ and $(\delta\Omega_0/\operatorname{Im}\Sigma_1)$. Recognizing that $\delta\Omega_0$ reverses sign when the tangential component $\mathbf{p}$ of the CE momentum is replaced by $-\mathbf{p}$, a successive-approximation method leads to the following chain of equations:

$$I_1\{\mathbf{f}_{\beta\alpha}^{(0)}(\mathbf{k},\omega)\}=0, \tag{24a}$$

$$\delta\Omega_0\mathbf{f}_{\beta\alpha}^{(0)}(\mathbf{k},\omega)=I_1\{\mathbf{f}_{\beta\alpha}^{(1)}(\mathbf{k},\omega)\}, \tag{24b}$$

$$(\omega-\Omega_0)\mathbf{f}_{\beta\alpha}^{(0)}(\mathbf{k},\omega)+\delta\Omega_0\mathbf{f}_{\beta\alpha}^{(1)}(\mathbf{k},\omega)$$
$$=\sigma_{\beta\alpha}+I_1\{\mathbf{f}_{\beta\alpha}^{(2)}(\mathbf{k},\omega)\}+I_2\{\mathbf{f}_{\beta\alpha}^{(0)}(\mathbf{k},\omega)\}, \tag{24c}$$

$$\mathbf{f}=\mathbf{f}^{(0)}+\mathbf{f}^{(1)}+\mathbf{f}^{(2)}+\ldots, \tag{24d}$$

where the collision integrals, which take into account the momentum and spin-orbit scattering $I_1\{\mathbf{f}\}$ and $I_2\{\mathbf{f}\}$, respectively, are given by

$$I_1\{f_{\beta\alpha}(\mathbf{k},\omega)\}=-2i\,\mathrm{Im}\,\Sigma_1(\mathbf{k},\mathbf{k};0-i\delta)f_{\beta\alpha}(\mathbf{k},\omega)+2i\pi\rho(\varepsilon_F)V\langle\mathcal{U}_1(\mathbf{k},\mathbf{k};\mathbf{k}_1,\mathbf{k}_1)f_{\beta\alpha}(\mathbf{k}_1,\omega)\rangle,$$
$$I_2\{f_{\beta\alpha}(\mathbf{k},\omega)\}=-2i\,\mathrm{Im}\,\Sigma_2(\alpha\mathbf{k},\alpha\mathbf{k},0-i\delta)f_{\beta\alpha}(\mathbf{k},\omega)+2i\pi\rho(\varepsilon_F)V\sum_{\gamma_1\gamma_2}\langle\mathcal{U}_2(\beta\mathbf{k},\alpha\mathbf{k};\gamma_1\mathbf{k}_1,\gamma_2\mathbf{k}_1)f_{\gamma_1\gamma_2}(\mathbf{k}_1,\omega)\rangle, \quad (25)$$

and we have omitted the indices of $\mathrm{Im}\,\Sigma_1$ and $\mathcal{U}_1$.

It follows from (24a) that $f^{(0)}(\mathbf{k},\omega)$ does not depend on the angles, i.e., on the direction of the CE momentum $\mathbf{k}$, this being a consequence of the general relation[10]

$$\mathrm{Im}\,\Sigma_1(\mathbf{k},\mathbf{k};0-i\delta)=\sum_{\mathbf{k}_1}\mathcal{U}_1(\mathbf{k},\mathbf{k};\mathbf{k}_1,\mathbf{k}_1)\,\mathrm{Im}\,\mathcal{G}(\mathbf{k}_1,\mathbf{k}_1;0-i\delta)=\pi\rho(\varepsilon_F)V\langle\mathcal{U}_1(\mathbf{k},\mathbf{k};\mathbf{k}_1,\mathbf{k}_1)\rangle, \quad (26)$$

where use was made of (22).

It follows from (14) and (15) that $\mathrm{Im}\,\Sigma_1$ and $\mathcal{U}_1$ do not depend on $\mathbf{p}$, therefore

$$f_{\beta\alpha}^{(1)}(\mathbf{k},\omega)=i\frac{\delta\Omega_0 f_{\beta\alpha}^{(0)}(\mathbf{k},\omega)}{2\,\mathrm{Im}\,\Sigma_1(\mathbf{k},\mathbf{k};0-i\delta)}=i\frac{\delta\Omega_0 f_{\beta\alpha}^{(0)}(\mathbf{k}\omega)}{\tau_V^{-1}+\tau_S^{-1}3\hat{q}^2}. \quad (27)$$

Substituting (27) in (24c) and averaging both halves of the result over the angles we obtain, taking (14), (15), (25), and (26) into account,

$$f_{\beta\alpha}^{(0)}(\mathbf{k},\omega)=A\sigma_{\beta\alpha}^{(n)}+B\sigma_{\beta\alpha}^{(t)}, \quad (28)$$

where

$$1/A=1/B+i/T_S,$$
$$1/B=\omega-\Omega_0+i/T_V+i\langle\delta\Omega_0^2/(\tau_V^{-1}+3\hat{q}^2\tau^{-1})\rangle+i/T_S, \quad (29)$$

while the normal and tangential components of the vector $\boldsymbol{\sigma}_{\beta\alpha}$ are given by

$$\sigma_{\beta\alpha}^{(n)}=\mathbf{n}(\boldsymbol{\sigma}_{\beta\alpha}\mathbf{n}),\quad \sigma_{\beta\alpha}^{(t)}=\boldsymbol{\sigma}_{\beta\alpha}-\mathbf{n}(\boldsymbol{\sigma}_{\beta\alpha}\mathbf{n}), \quad (30)$$

where $\mathbf{n}$ is the normal to the surface.

We note that $A\neq B$ in (29). This is due to the anisotropy of the spin-orbit interaction of the CE with the surface centers:

$$\sum_{\gamma_1\gamma_2}\int\frac{d\Omega\,d\Omega_1}{(2\pi)^2}\mathcal{U}_2{}^s(\beta\mathbf{k},\alpha\mathbf{k};\gamma_1\mathbf{k}_1,\gamma_2\mathbf{k}_1)\sigma_{\gamma_1\gamma_2}^{(t)}=0.$$

Since the Z axis was directed in spin space along the static field $\mathbf{H}_0$ and the X axis along the RF field $\mathbf{H}_1$, it follows that the Y axis is directed along $\mathbf{H}_0\times\mathbf{H}_1$, so that

$$\sigma_{\beta\alpha}^{(n)}\mathbf{H}_1(\omega))=(\cos^2\widehat{\mathbf{nH}_1}-i\cos\widehat{\mathbf{nH}_1}\cos\widehat{\mathbf{n}[\mathbf{H}_0\times\mathbf{H}_1]})H_1(\omega)$$
$$(\sigma_{\beta\alpha}^{(t)}\mathbf{H}_1(\omega))=(\sin^2\widehat{\mathbf{nH}_1}+i\cos\widehat{\mathbf{nH}_1}\cos\widehat{\mathbf{n}[\mathbf{H}_0\times\mathbf{H}_1]})H_1(\omega), \quad (31)$$

where we have used the fact that $\sigma_{\beta\alpha}^X=1$ and $\sigma_{\beta\alpha}^Y=-i$.

Substituting (28) in (21), we obtain, taking (29) and (31) into account,

$$M^-(\omega)=V\chi_0\left\{1-\frac{\omega\cos^2\widehat{\mathbf{nH}_1}}{\omega-\Omega_0+iT_n^{-1}}-\frac{\omega\sin^2\widehat{\mathbf{nH}_1}}{\omega-\Omega_0+iT_t^{-1}}+\frac{\omega(T_n^{-1}-T_t^{-1})\cos\widehat{\mathbf{nH}_1}\cos\widehat{\mathbf{n}[\mathbf{H}_0\times\mathbf{H}_1]}}{(\omega-\Omega_0+iT_n^{-1})(\omega-\Omega_0+iT_t^{-1})}\right\}H_1(\omega), \quad (32)$$

and with allowance for (9) and (10) we obtain for the RF-field power absorbed in the film the expression

$$P=\frac{\omega^2H_1^2V}{4}\chi_0\left\{\frac{T_n^{-1}\cos^2\widehat{\mathbf{nH}_1}}{(\omega-\Omega_0)^2+T_n^{-2}}+\frac{T_t^{-1}\sin^2\widehat{\mathbf{nH}_1}}{(\omega-\Omega_0)^2+T_t^{-2}}+\frac{(\Omega_0-\omega)(T_n^{-2}-T_t^{-2})\cos\widehat{\mathbf{nH}_1}\cos\widehat{\mathbf{n}[\mathbf{H}_0\times\mathbf{H}_1]}}{[(\omega-\Omega_0)^2+T_n^{-2}][(\omega-\Omega_0)^2+T_t^{-2}]}\right\}. \quad (33a)$$

In (32) and (33) we have introduced the notation

$$T_n^{-1}=T_V^{-1}+(T_S^n)^{-1}+(T_S')^{-1},$$
$$T_t^{-1}=T_V^{-1}+(T_S^t)^{-1}+(T_S')^{-1}, \quad (34)$$

where T_S^n and T_S^t are given by

$$(T_S^n)^{-1}=2T_S^{-1},\ (T_S^t)^{-1}=T_S^{-1}, \quad (35a)$$

and the additional contribution to the total line width $(T_S')^{-1}$ is given by

$$(T_S')^{-1}=\begin{cases}\langle\delta\Omega_0^2\rangle\tau_V, & \tau_V^{-1}\gg3\tau_S^{-1},\\ \frac{7}{9}\langle\delta\Omega_0^2\rangle\tau_S, & \tau_V^{-1}\ll3\tau_S^{-1},\end{cases} \quad (36a)$$
$$\langle\delta\Omega_0^2\rangle=\sin^2\mathbf{nH}_0\frac{c_S^2}{d^2}\frac{b_S^2}{\hbar^2}\frac{64}{35}.$$

We can write $(T_S^n)^{-1}$ and $(T_S^t)^{-1}$ in the Dyson form (2), if we define ε as follows:

$$\varepsilon=\langle\mathbf{jS}\rangle^{-1}\left\langle\sum_{\mathbf{k}_1}W(\alpha\mathbf{k}_1,\beta\mathbf{k})\right\rangle, \quad (37)$$

where $W(1,2)$ is the probability of the transition of the CE per unit time from the state $|2\rangle$ into the state $|1\rangle$, $\langle\ldots\rangle$ denotes averaging over the angles, and $\mathbf{j}$ is the probability flux density of the CE on the surface $\mathbf{S}=\mathbf{n}S$. It follows from (4) that $\langle\mathbf{jS}\rangle=v_F/4d$. Calculating $W(1,2)$ in the Born approximation with the Hamiltonian $\mathcal{H}_2^S$ given in (8b), we obtain

$$\varepsilon=\frac{2}{v_F}c_S\frac{2\pi}{\hbar}\rho(\varepsilon_F)b_S^2\frac{16}{45}(1+\cos^2\widehat{\mathbf{nH}_1}), \quad (38)$$

where we have directed the quantization axis along $\mathbf{H}_1$, since we are considering the relaxation of the magnetization component $\mathbf{M}\parallel\mathbf{H}_1$, which is responsible for the absorption of the RF power in the film. Taking (16), (35), and (38) into account we obtain

$$(T_S^n)^{-1}=\varepsilon_n\frac{v_F}{2d},\quad (T_S^t)^{-1}=\varepsilon_t\frac{v_F}{2d}, \quad (35b)$$

where ε_n and ε_t are determined from (38) for the angles $\widehat{\mathbf{nH}_1}=0$ and $\widehat{\mathbf{nH}_1}=\pi/2$ respectively. The quantities ε_n and ε_t have the meaning of the probabilities of the spin flip of the CE in a single collision with the surface $z=0$, if the spin is normal and tangential to the surface, respectively.

Expression (36) for $(T_S')^{-1}$ can be rewritten, taking (38) into account, in the form

$$(T_S')^{-1}=\begin{cases}\frac{9}{7}\theta\varepsilon_t\frac{v_Fl_V}{d^2}\sin^2\widehat{\mathbf{nH}_0}, & l_V\ll l_S=\frac{d}{\eta},\\ \theta\frac{\varepsilon_t}{\eta}\frac{v_F}{2d}\sin^2\widehat{\mathbf{nH}_0}, & l_V\gg l_S=\frac{d}{\eta},\end{cases} \quad (36b)$$

where $\theta = 4\pi c_S k_F^2$ is a quantity proportional to the degree of coverage of the surface $z = 0$ by the scattering centers, while $l_V = \tau_V v_F$ and $l_S = \tau_S v_F$ are the CE mean free paths connected with the volume and surface momentum scattering, with l_S expressed in terms of the thickness of the film d and the diffuseness coefficient η of the surface $z = 0$ for electrons traveling in a direction normal to the surface.[12]

Since the quantity observed in PRCE experiments is not the power P itself but the derivative dP/dH_0, it follows that the observed PRCE signal is described by the expression

$$\frac{dP}{dH_0} = C\left\{-\frac{2\varkappa\alpha\cos^2\widehat{\mathbf{nH}}_1}{(\alpha^2+\varkappa^2)^2} - \frac{2\alpha\sin^2\widehat{\mathbf{nH}}_1}{(\alpha^2+1)^2} + \frac{[\varkappa^2-\alpha^2(\varkappa^2+1)-3\alpha^4](\varkappa^2+1)\cos\widehat{\mathbf{nH}}_1\cos\mathbf{n}\widehat{[\mathbf{H}_0\mathbf{H}_1]}}{(\alpha^2+\varkappa^2)^2(\alpha^2+1)^2}\right\}. \quad (33b)$$

Here

$$C = \frac{\omega^2 H_1^2 V}{4}\chi_0\mu T_t^2, \quad \alpha = (\Omega_0-\omega)T_t, \quad \varkappa = \frac{T_n^{-1}}{T_t^{-1}}.$$

4. DISCUSSION OF THE RESULTS

If the angle $\widehat{\mathbf{nH}}_1$ is equal to 0 or $\pi/2$, then the PRCE signal, as follows from (33) and (34), takes the form of a Lorentz absorption line with total width T_n^{-1} or T_t^{-1}, consisting of a volume contribution T_V^{-1}, a Dyson contribution $(T_S^n)^{-1}$ or $(T_S^t)^{-1}$, and an additional surface contribution $(T_S')^{-1}$.

As seen from (36b), $(T_S')^{-1}$ increases with increasing mean free path and flattens out at $l_V \gg d/\eta$, where it assumes a maximum value that depends on the surface diffuseness coefficient η.

The appearance of $(T_S')^{-1}$ in the total line width is due to the fact that the impurities located on only one surface produce in the film a nonzero average electric field $E \sim d^{-1}(\mathbf{E} \parallel \mathbf{n})$. Therefore the CE spin moving with momentum $\mathbf{k}$ in this electric field is acted upon by a magnetic field $\mathbf{B}(\mathbf{k}) \sim \mathbf{E}\times\mathbf{k}$ $(\mathbf{B}(\mathbf{k}) \perp \mathbf{n})$, which changes the frequency of the CE spin precession about the static magnetic field $\mathbf{H}_0$ by an amount $\delta\Omega_0 \sim |\mathbf{B}(\mathbf{k})| \sin\widehat{\mathbf{nH}}_0 \sim d^{-1} \times \sin\widehat{\mathbf{nH}}_0$. The field $\mathbf{B}(\mathbf{k})$ changes its direction randomly within the electron mean free path time τ, and this leads[13] to an additional contribution to the line width, equal to $\langle \delta\Omega_0^2 \rangle \tau$ and of the same order of magnitude as (36a). If identical impurities are randomly disposed on both surfaces of the film, then the average electric field produced by them in the film is $\mathbf{E} = 0$, so that $(T_S')^{-1} = 0$.

Let us estimate the maximum value of $(T_S')^{-1}$ for Li films. At $c_S = 10^{14}$ cm^{-1} we have $(T_S')^{-1} = 0.21\eta^{-1}(T_S^t)^{-1}$. It follows from (17) that it is necessary to choose $\eta < 1$. It is seen that $(T_S')^{-1}$ can be larger than the Dyson contribution (35).

From the experimental values of the two surface contributions $(T_S^t)^{-1}$ and $(T_S')^{-1}$ we can determine, by using (35) and (36), the values of ε_t and of the concentration c_S and consequently, by using (38), also the constant b_S of the interaction of the CE spins with the surface centers.

We note that the mean free path l_V in (36b) can be taken to mean the volume mean free path corresponding to the given temperature T, if T is higher than the Debye temperature T_D, for in this case the scattering of the CE by the phonons in the metal is isotropic, and consequently is equivalent to the scattering in the model of short-range centers. It follows therefore from (36b) that the total line width increases with decreasing temperature. It was noted above that the surface-temperature contribution to the line width $(T_S')^{-1} = 0$ if the scattering centers are symmetrically distributed over both surfaces, as is expected in colloids of lithium in LiF crystals. It follows from the results of this study that the observed increase of the total width of the PRCE line[3,4] with decreasing temperature in radiation colloids of lithium is not due to surface relaxation. The difference between (36) and (3) is due to the fact that in[3,4] the frequency of the collisions of the CE with the surface has been assumed to be proportional to the mean free path l_V, although it is known[14] that this assumption is incorrect.

We note that the definition (37) of ε differs from the expression obtained earlier[5] from an analysis of the scattering of a wave packet by a surface. In (37), $(\mathbf{j}\cdot\mathbf{S})$ and $W(1, 2)$ are averaged over the angles separately, while in[5] it is the quotient $W(1, 2)/(\mathbf{j}\cdot\mathbf{S})$ which is averaged. Expression (37) must be regarded as a correct definition of ε, since the more effective momentum scattering averages the probability flux density $(\mathbf{j}\cdot\mathbf{S})$ incident on the surface, i.e., it averages the frequency of the CE collisions with the surface.

If the angle between $\mathbf{n}$ and $\mathbf{H}_1$ is arbitrary, then it is seen from (33) that the line shape contains a "dispersion" term that leads to asymmetry of the observed line dP/dH_0 even in very thin metallic plates [Eq. (5)].

The appearance of this term can be understood in the following manner. The RF-field power absorbed by the sample is determined in the rotating coordinate frame[8] by the magnetization component M_y. Owing to the anisotropy of ε, the magnetization component M_x, which is initially directed along the x axis in the rotating coordinate system, will no longer be exactly directed along x, and a y projection will appear, $\Delta M_y \sim M_x^0(T_n^{-1} - T_t^{-1})$. Since $M_x^0 \sim (\omega - \Omega_0)$, a term $\sim (\omega - \Omega_0)(T_n^{-1} - T_t^{-1})$ will appear in the line shape.

The asymmetry of the PRCE line is characterized[15] by the ratio C/D of the shoulders of the observed signal. It follows from (33b) that C/D is a function of the surface-relaxation anisotropy coefficient $\varkappa = T_n^{-1}/T_t^{-1}$. The values of C/D for different values of $\varkappa$ and different film orientations are shown in Fig. 2. In the limiting case when the anisotropy of the surface relaxation is substantial, $l_V \ll d$ and $T_V^{-1} \ll (T_S^t)^{-1}$, the quantity $\varkappa$ assumes a maximum value $\varkappa_{max} = \varepsilon_n/\varepsilon_t$. In the assumed model $\varkappa_{max} = 2$, as follows from (38). If $\varkappa = 2$ and the angle between $\mathbf{n}$ and $\mathbf{H}_1$ is equal to the angle between $\mathbf{n}$ and $\mathbf{H}_0\times\mathbf{H}_1$, which in turn is equal to $\pi/4$, then $C/D = 1.71$. For thicknesses $d \geq 0.5\delta$ and for an angle $\pi/2$ between $\mathbf{n}$ and $\mathbf{H}_1$, the experimental value of the coefficient C/D of the asymmetry due to the skin effect were obtained experimentally in[15] namely $C/D = 1.85$ for d

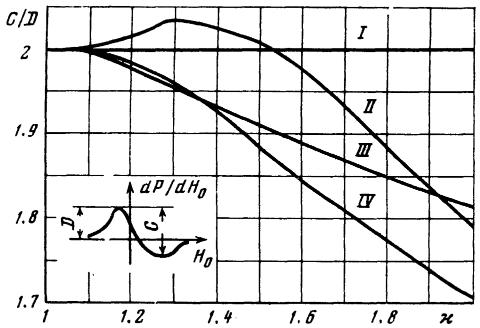

FIG. 2. Dependence of the asymmetry parameter C/D of the absorption line on the ratio $\varkappa = T_n^{-1}/T_t^{-1}$ in the case H_0 n. The curves correspond to the following film orientations: I $-\widehat{nH_1} = 0$ or $\pi/2$; II—$\widehat{nH_1} = \pi/8$; III—$\widehat{nH_1} = 3\pi/8$; IV—$\widehat{nH_1} = \pi/4$.

$= 0.5\delta$, which is larger than the estimate $(C/D)_{min} = 1.71$ given above. It is therefore necessary to take into account the asymmetry introduced into dP/DH_0 by the surface relaxation when the film thickness as measured by the PRCE method,[15] for otherwise the experimental values can exceed the real thicknesses by an order of magnitude.

In addition, the dependence of the line shape on the sample orientation can be used to determine the shapes of the radiation colloids, since the orientation dependence of the observed PRCE signals in spherical colloids is expected to be weaker than in colloids in the form of films.

We note in conclusion that expression (32) for the magnetization can be obtained from a solution of the Bloch equation with a modified boundary condition (1)

$$2\frac{\varepsilon_n^{\text{eff}}}{L}\mathbf{M}_n(\mathbf{r},t) = -\left(\mathbf{n}\frac{\partial}{\partial\mathbf{r}}\right)\mathbf{M}_n(\mathbf{r},t), \quad z=0,$$
$$2\frac{\varepsilon_t^{\text{eff}}}{L}\mathbf{M}_t(\mathbf{r},t) = -\left(\mathbf{n}\frac{\partial}{\partial\mathbf{r}}\right)\mathbf{M}_t(\mathbf{r},t), \quad z=0, \tag{39}$$

where $\mathbf{M}_n = \mathbf{n}(\mathbf{M}\cdot\mathbf{n})$ and $\mathbf{M}_t = \mathbf{M} - \mathbf{n}(\mathbf{M}\cdot\mathbf{n})$ are the normal and tangential components of the magnetization, $L = 4D/v_F$ as in[1,2],

$$\varepsilon_n^{\text{eff}} = \varepsilon_n + \sin^2\widehat{\mathbf{nH}_0}\frac{\beta}{d},$$
$$\varepsilon_t^{\text{eff}} = \varepsilon_t + \sin^2\widehat{\mathbf{nH}_0}\frac{\beta}{d}, \tag{40}$$

where ε_n and ε_t are the probabilities of the CE spin flip in collisions with the surface for a spin direction normal and tangential to the surface, $\beta = \frac{9}{7}\theta\varepsilon_t l_V$ at $l_V \ll d/\eta$ and $\beta = \theta\varepsilon_t d/\eta$ at $l_V \gg d/\eta$ and $\theta = 4\pi c_S/k_F^2$.

Expression (32) can also be obtained by solving the diffusion equation, by replacing in formula (20) of[1] the product $(\mathbf{sH}_1)\mathcal{G}(\mathbf{r}', t'; \mathbf{r}, t)$ by $(\mathbf{s}_n\mathbf{H}_1)\mathcal{G}_n(\mathbf{r}', t'; \mathbf{r}, t) + (\mathbf{s}_t\mathbf{H}_1) \times \mathcal{G}_t(\mathbf{r}', t'; \mathbf{r}, t)$, where $\mathcal{G}_n(1, 2)$ and $\mathcal{G}_t(1, 2)$ satisfy boundary conditions analgous to (39), $\mathbf{s}_n = \mathbf{n}(\mathbf{s}\cdot\mathbf{n})$, and $\mathbf{s}_t = \mathbf{s} - \mathbf{n}(\mathbf{s}\cdot\mathbf{n})$. The function $\mathcal{G}_n(\mathbf{r}', t'; \mathbf{r}, t)$ has the meaning of the probability of observing at the point r at the instant t an electron initially located at $\mathbf{r}'$ at the instant t' with a spin $\mathbf{s}$ directed normal to the surface. The interpretation of $\mathcal{G}_t(\mathbf{r}', t'; \mathbf{r}, t)$ is analogous.

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