SPIN ECHO IN NONFERROMAGNETIC METALS

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A theory of the linearized spin echo in a degenerate electron liquid is developed. The case is considered of arbitrary orientation of the wave vector of the spin wave and of the gradient of the magnetic field relative to the direction of the external magnetic field. It is shown that it is possible to determine the imaginary part of the generalized diffusion coefficient of the spin wave from experiments on spin echo.

IT is known that spin waves can be excited in a charged Fermi liquid placed in an external homogeneous magnetic field[1]. These spin waves were recently observed experimentally and some information was obtained concerning the parameters of the Fermi liquid. In[2] it was shown that if spin-echo experiments are performed in a neutral Fermi liquid not in a constant magnetic field but in a field with a small spatial gradient, then such experiments can yield additional information on the interaction parameters of the Fermi liquid. Experiments of this type are preferable to the first experiments performed on the observation of spin waves in alkali metals[3], since they are not connected with the boundary conditions on the surface of the metal or of the neutral liquid, and their theory is relatively simple. A neutral Fermi liquid (He^3 or its solution in He^4) were chosen for the echo experiments because of the absence of a coupling between the spin waves and the orbital motion. The results of a possible echo experiment, also on a neutral Fermi liquid in a magnetic field with a small gradient along the magnetic-field force line, were expressed in[4] in terms of the parameters of spin waves propagating in the same direction.

We consider in the present paper a similar problem for a degenerate electronic Fermi liquid. We carry out the analysis for arbitrary directions of the gradient of the magnetic field and of the wave vector of the spin wave relative to the magnetic field.

The gist of the spin-echo experiment consists in the following. In an external constant magnetic field, an equilibrium distribution of magnetization is established along this field. At the initial instant of time, a high-frequency pulse is used to excite a definite mode of transverse spin density. We confine ourselves to the linear approximation for the transverse spin density, i.e., to HF pulses at small angles to the direction of the magnetic field. In a weakly inhomogeneous magnetic field, an irreversible decay of the transverse spin density will occur, first because of the collisions with the impurities and, second, because of precession of the transverse moment in the inhomogeneous magnetic field. The reversible dephasing of the spin density, produced by the magnetic field, is separated from the irreversible decay by using a subsequent 180° HF pulse at the instant t_0, which reverses the sign of the projection of the total angular momentum on the direction of the magnetic field and produces the echo effect at the instant $2t_0$.

To describe the time behavior of the spin density we use the kinetic equations obtained in[1]. In the linear approximation, the oscillations of the phase spin density are practically unrelated to the oscillations of the quasiparticle distribution (a view of the smallness of the ratio of the energy of the electron level splitting in the magnetic field to the Fermi energy). We also neglect the small forces that result from the gradient of $\mathbf{B}(\mathbf{r})$ and lead to a change in the equilibrium distribution of the quasiparticles; these gradients will be taken into account only in the determination of the time development of the modes of the phase spin density. Therefore, at the initial instant the equilibrium spin density is determined by the expression

$$\sigma_0 = -\frac{\mu}{1+\beta_0}\frac{\partial f_0}{\partial \varepsilon}\mathbf{B}(\mathbf{r}) = -\gamma\frac{\partial f_0}{\partial \varepsilon}\mathbf{B}(\mathbf{r}), \qquad (1)$$

where $f_0(\mathbf{p})$ is the Fermi distribution function, β_0 is the first coefficient of the expansion of the quasiparticle interaction function, which depends on the spins,

$$\Psi(\mathbf{p},\mathbf{p}') = \frac{\pi^2\hbar^3 v_F}{p_F^2}\sum_{l=0}^{\infty}(2l+1)\beta_l P_l(\cos\chi),$$

where χ is the angle between the vectors $\mathbf{p}$ and $\mathbf{p}'$, and it is assumed that the Fermi surfaces differ little from spherical, a condition well satisfied for alkali metals.

The linearized equation for the spin-density component transverse to the magnetic field can be written in the form

$$\frac{\partial\sigma_\perp}{\partial t} + \left(\mathbf{v}_F\frac{\partial}{\partial\mathbf{r}}\right)\left(\sigma_\perp - \frac{\partial f_0}{\partial\varepsilon}\int d\mathbf{p}'\,\Psi(\mathbf{p},\mathbf{p}')\sigma_\perp(\mathbf{p}')\right) - \frac{2\gamma}{\hbar}\left[\mathbf{B},\left(\sigma_\perp - \frac{\partial f_0}{\partial\varepsilon}\int d\mathbf{p}'\,\Psi\sigma_\perp\right)\right] + \frac{e}{c}[\mathbf{v}_F\mathbf{B}]\left(\frac{\partial\sigma_\perp}{\partial\mathbf{p}} - \frac{\partial f_0}{\partial\varepsilon}\int d\mathbf{p}'\,\Psi\frac{\partial\sigma_\perp}{\partial\mathbf{p}'}\right) = \left(\frac{\partial\sigma_\perp}{\partial t}\right)_{\text{St}}. \qquad (2)^*$$

In the case of low temperatures, the collision integral is given by

$$\left(\frac{\partial\sigma_\perp}{\partial t}\right)_{\text{St}} = -\left(\frac{1}{\tau_0}+\frac{1}{T}\right)\left(\sigma_\perp - \frac{\partial f_0}{\partial\varepsilon}\int d\mathbf{p}'\,\Psi\sigma_\perp\right) + \frac{\partial f_0}{\partial\varepsilon}\left(\int d\mathbf{p}'\frac{\partial f_0}{\partial\varepsilon'}\right)^{-1}\int\frac{d\mathbf{p}'}{\tau(\chi)}\left(\sigma_\perp - \frac{\partial f_0}{\partial\varepsilon'}\int d\mathbf{p}''\,\Psi\sigma_\perp\right),$$

where T is the spin-flip time, and τ_0 and $\tau(\chi)$ are parameters characterizing the relaxation of the quasi-

*$[\mathbf{v}_F\mathbf{B}] \equiv \mathbf{v}_F \times \mathbf{B}$.

particle momentum. For the function $\tau(\chi)$ we shall use the expansion

$$\frac{1}{\tau(\chi)}=\sum_{l=0}^{\infty}\frac{(2l+1)}{\tau_l}P_l(\cos\chi).$$

We have left out from (2) terms containing the magnetic field of the perturbation. Such an approximation is justified because of the very small static paramagnetic susceptibility of the electron gas ($\chi_0 = \mu^2 p_F^2/\pi^2 h^3 v_F$) and leads to the same spin-wave spectrum which is obtained from the pole of the HF magnetic susceptibility if the indicated terms are retained in Eq. (2) and this equation is supplemented with Maxwell's equations.

Expanding the transverse spin density in terms of the angle components $\sigma^{\pm} = \sigma_x \pm i\sigma_y$, we can obtain from the system (2) two independent equations for each of the components. It is convenient to introduce the notation

$$\sigma^{\pm}=\frac{\partial f_0}{\partial \varepsilon}m^{\pm}(\mathbf{p},\mathbf{r},t),\quad \overline{\Psi}=\frac{2p_F^2}{(2\pi\hbar)^3 v_F}\Psi.$$

Choosing then a weakly inhomogeneous field in the form $\mathbf{B}(\mathbf{r}) = (1 + \mathbf{k}\cdot\mathbf{r})B_0$ and changing over to Fourier components

$$m^{\pm}(\mathbf{r})=\int \exp(i\mathbf{q}\mathbf{r})\,m^{\pm}(\mathbf{q})\,d\mathbf{q},$$

we obtain for the quantities $m^{\pm}(\mathbf{q})$ the equations

$$\frac{\partial}{\partial t}m^{\pm}(\mathbf{q},\mathbf{p},t)+\left(iv_F\mathbf{q}-\Omega\frac{\partial}{\partial\varphi}\mp i\Omega_0+\frac{1}{\tau_0}+\frac{1}{T}\right)\left(m^{\pm}+\int d\Omega'\overline{\Psi}m^{\pm}\right)$$
$$+\left(\mathbf{k}\frac{\partial}{\partial\mathbf{q}}\right)\left[\left(\pm\Omega_0-i\Omega\frac{\partial}{\partial\varphi}\right)\left(m^{\pm}+\int d\Omega'\overline{\Psi}m^{\pm}\right)\right]=$$
$$=\int\frac{d\Omega'}{4\pi\tau(\chi)}\left(m^{\pm}+\int d\Omega''\overline{\Psi}m^{\pm}\right), \qquad (3)$$

where $\Omega = eB_0/mc$ is the cyclotron frequency of the electron, $\Omega_0 = 2\gamma B/\hbar$, φ is the azimuthal angle in momentum space, where the polar axis is oriented along the magnetic field (the z axis), and $d\Omega$ is the solid-angle element.

We shall find it convenient to obtain solutions of (3) in the case of a homogeneous magnetic field ($\mathbf{k} = 0$). These solutions are simplest to obtain for long-wave oscillations ($\mathbf{q}\cdot\mathbf{v}_F \ll \Omega$), by using perturbation theory. In the zeroth approximation ($\mathbf{q} = 0$) the eigenfunctions of the spin density are the spherical harmonics $Y_{lm}(\theta,\varphi)$, and the eigenfrequencies are given by the formula

$$\omega_{lm}^{\pm}(0) = (1+\beta_l)(m\Omega\pm\Omega_0)+i\nu_l,\ \mathbf{q}=0, \qquad (4)$$

where $\nu_l = (1 + \beta_l)(1/\tau_0 + 1/T - 1/\tau_l)$. We note that the time dependence of the perturbations of the spin density was sought in the form $\sim\exp(i\omega_{lm}t)$. In the next order in q we obtain for the eigenfunctions and frequencies of the spin waves in a homogeneous magnetic field the expressions

$$\mu_{lm}^{\pm}(\mathbf{q}) = Y_{lm}(\theta,\varphi) - v_F q\left[\frac{4\pi}{3(2l+1)}\right]^{1/2}\cdot$$
$$\times\sum_{l'm'}\frac{(2l'+1)^{1/2}(1+\beta_l)}{\omega_{lm}^{\pm}(0)-\omega_{l'm'}^{\pm}(0)}C_{l'0;\,10}^{l0}C_{l'm';\,1m-m'}^{lm}Y_{1,m-m'}(\theta_1,\varphi_1)Y_{l'm'}(\theta,\varphi),$$
$$\omega_{lm}^{\pm}(\mathbf{q}) = \omega_{lm}^{\pm}(0) + q^2 D_{lm}^{\pm}(\mathbf{q}), \qquad (5)$$

where the generalized diffusion coefficient is given by

$$D_{lm}^{\pm}(\mathbf{q}) = \frac{4\pi v_F^2}{3(2l+1)}(1+\beta_l)\cdot$$
$$\times\sum_{l'm'}\frac{(1+\beta_{l'})(2l'+1)}{\omega_{lm}^{\pm}(0)-\omega_{l'm'}^{\pm}(0)}\left|C_{l'0;\,10}^{l0}C_{l'm';\,1m-m'}^{lm}Y_{1,m'-m}(\theta_1,\varphi_1)\right|^2. \qquad (6)$$

Here θ_1 and φ_1 are the polar angles of the wave vector q, and $C^{\cdot\cdot}_{\cdots}$ are the Clebsch-Gordan coefficients. It is seen from (5) that mixing of the modes with different values of l takes place when $q \neq 0$. Particular interest attaches to the mode $l = m = 0$, since it makes a main contribution to the macroscopic magnetization density ($\mathbf{M} = \int\sigma\cdot d\mathbf{p}$). For the mode $l = 0$ we have

$$D_{00}(\mathbf{q})=\frac{v_F^2}{3}(1+\beta_0)(1+\beta_1)\left[\frac{\cos^2\theta_1}{\omega_{00}-\omega_{10}}\right.$$
$$\left.+\sin^2\theta_1\frac{\omega_{00}-\omega_{10}}{(\omega_{00}-\omega_{10})^2-\Omega^2(1+\beta_1)^2}\right],$$
$$\mu_{00}(\mathbf{q})=\frac{1}{\sqrt{4\pi}}+\frac{v_F q(1+\beta_0)}{\sqrt{4\pi}}\left\{-\frac{\cos\theta_1\cos\theta}{\omega_{00}-\omega_{10}}\right.$$
$$+\frac{\sin\theta_1\sin\theta}{(\omega_{00}-\omega_{10})^2-\Omega^2(1+\beta_1)^2}\left[(\omega_{00}-\omega_{10})\cos(\varphi_1-\varphi)\right.$$
$$\left.\left.-i\Omega(1+\beta_1)\sin(\varphi_1-\varphi)\right]\right\}. \qquad (7)$$

In a weakly inhomogeneous magnetic field, it is convenient to change over in (3) to the interaction representation and to seek the solutions in the form

$$m^{\pm}(\mathbf{q},\mathbf{p},t)=\sum_{lm}\Phi_{lm}^{\pm}(\mathbf{q},t)\,\mu_{lm}^{\pm}(\mathbf{q},\Omega)\exp\left[i\omega_{lm}^{*\pm}(\mathbf{q})t\right], \qquad (8)$$

where μ_{lm} and ω_{lm} are the eigenfunctions and frequencies of Eq. (3) at $k = 0$. Substituting (8) in (3), multiplying it by μ_{lm}, and integrating over the angles Ω, we obtain the following system of equations for the determination of the functions Φ_{lm}:

$$\frac{\partial\Phi_{lm}^{\pm}}{\partial t}+\alpha_{lm}^{\pm}\left(\mathbf{k}\frac{\partial\Phi_{lm}^{\pm}}{\partial\mathbf{q}}+it\Phi_{lm}^{\pm}\mathbf{k}\frac{\partial\omega_{lm}^{\pm}}{\partial\mathbf{q}}\right)$$
$$+\sum_{l'm'}\exp(-i\Delta\omega t)\frac{v_F}{\Delta\omega}\left[\frac{4\pi(2l'+1)}{3(2l+1)}\right]^{1/2}C_{l'0;\,10}^{l0}C_{l'm';\,1m-m'}^{lm}$$
$$\times\left\{k(1+\beta_{l'})\alpha_{lm}^{\pm}\Phi_{l'm'}^{\pm}Y_{1,m'-m}(\theta_0,\varphi_0)-\frac{q}{\Delta\omega^*}Y_{1m'-m}(\theta_1,\varphi_1)\right.$$
$$\times\left[\frac{\partial\Phi_{l'm'}^{\pm}}{\partial t}\langle(\beta_l-\beta_{l'})\,\mathrm{Re}\,\Delta\omega+i(2+\beta_l+\beta_{l'})\,\mathrm{Im}\,\Delta\omega\rangle\right.$$
$$+(1+\beta_l)(1+\beta_{l'})\left(\mathbf{k}\frac{\partial\Phi_{l'm'}^{\pm}}{\partial\mathbf{q}}+it\Phi_{l'm'}^{\pm}\mathbf{k}\frac{\partial\omega_{l'm'}^{\pm}}{\partial\mathbf{q}}\right)$$
$$\left.\left.\times\langle(m'-m)\Omega\,\mathrm{Re}\,\Delta\omega+i(\pm2\Omega_0+m'\Omega+m\Omega)\,\mathrm{Im}\,\Delta\omega\rangle\right]\right\}=0, \qquad (9)$$

where $\alpha_{lm}^{\pm} = (\pm\Omega_0 + m\Omega)(1 + \beta_l)$, $\Delta\omega = \omega_{lm}(0) - \omega_{l'm'}(0)$. In deriving this system we used the approximation linear in q for the eigenfunctions μ_{lm} (see Eq. (5)); the contribution from the terms of the eigenfunctions containing higher powers of q turns out to be negligible. It should be noted that the sum in (9) does not contain a term with $l' = l$.

We are interested in asymptotic solutions of the system (9) for $t \gg \Omega^{-1}$. At $k = 0$, the functions Φ_{lm} do not depend on the time, and their values represent the initial conditions of the problem, characterizing the state of the system prior to turning on the interaction ($k \neq 0$ at $t > 0$). Let us assume that at the instant $t = 0$ only one of the functions differs from zero, $\Phi_{lm}(\mathbf{q}, t = 0) = m_{l_0m_0}(\mathbf{q})\,\delta_{ll_0}\delta_{mm_0}$, where $m_{l_0m_0}$ is the Fourier coefficient of the spin density of the mode l_0m_0 at the instant $t = 0$, i.e., the HF pulse excites only one transverse spin density mode. In this case the contribution of the nondiagonal terms in the equation for the function $\Phi_{l_0m_0}$ can be neglected and we obtain for it the following equation (the subscript zero will henceforth be omitted wherever this can cause no confusion):

$$\frac{\partial \Phi_{lm}}{\partial t} + \alpha_{lm}\left(\mathbf{k}\frac{\partial \Phi_{lm}}{\partial \mathbf{q}} + it\mathbf{k}\frac{\partial \omega_{lm}}{\partial \mathbf{q}}\Phi_{lm}\right) = 0.$$

Integrating this equation among the characteristic $\mathbf{c} = \mathbf{q} - \alpha_{lm}t\mathbf{k}$, we get a solution in the form

$$\Phi_{lm}(\mathbf{q}, t) = m_{lm}(\mathbf{c}) \exp\left[-i\alpha_{lm}\int_0^t t'\mathbf{k}\frac{\partial}{\partial \mathbf{c}}\omega_{lm}(\mathbf{c} + \alpha_{lm}t'\mathbf{k})\,dt'\right]$$

$$= m_{lm}(\mathbf{c}) \exp\left[-i\alpha_{lm}\frac{t^2}{2}\mathbf{k}\frac{\partial \omega_{lm}(\mathbf{c})}{\partial \mathbf{c}} - i\alpha_{lm}^2\frac{t^3}{3}\left(\mathbf{k}\frac{\partial}{\partial \mathbf{c}}\right)^2\omega_{lm}(\mathbf{c}) - \ldots\right]. \tag{10}$$

The functions Φ_{lm}, which were equal to zero at the initial instant of time, are determined from the system (9) by equating the derivative $\partial \Phi_{lm}/\partial t$ to the sum of the nondiagonal terms, in which only one term containing $\Phi_{l_0 m_0}$ is retained. It is easy to verify that when $t \gg \Omega^{-1}$ their order of magnitude is

$$\Phi_{lm} \sim \left[\alpha_{l_0 m_0}\frac{\mathbf{k}\mathbf{v}_F}{|\Delta\omega|^2}\right]^{|l-l_0|}\Phi_{l_0 m_0}. \tag{11}$$

By assumption, the parameter $\mathbf{k}\cdot\mathbf{v}_F/\Delta\omega \ll 1$, and therefore the nondiagonal modes $(l - l_0)$ turn out to be small. We see now that the neglect of the nondiagonal terms (the sum) in the equation for the function $\Phi_{l_0 m_0}$ was justified, since these terms result in a correction of the order of

$$\delta\Phi_{l_0 m_0} \sim \alpha^2_{l_0 m_0}\frac{(\mathbf{k}\mathbf{v}_F)^2}{|\Delta\omega|^4}\Phi_{l_0 m_0},$$

which is negligibly small for the asymptotic solution (10), smaller by a factor $\Omega^3 t^3$ than the terms taken into account by this solution.

We note that in formula (10) there were retained the first two terms of the expansion in powers of k, since in the long-wave limit (q = 0) the terms proportional to the odd powers of k vanish. In this limit, the ratio of the next effective term to that taken into account by us is of the order of

$$k^2\alpha_{lm}^2 t^2\omega^{\mathrm{IV}}/\omega'' \sim k^2 v_F^2 t^2 < 1.$$

Thus, if we expand in powers of k in (10) and discard terms with higher powers of k, then a limitation results on the observation time. However, for weakly-inhomogeneous magnetic fields this limitation is weak and the time of observation can exceed by many orders of magnitude the Zeeman period (Ω^{-1}).

We now return to the spin-echo experiment. As already noted earlier, if we excite the mode m^+_{lm} at the initial instant, then in the succeeding instant t_0 the external HF signal transforms it into the mode $m^-_{l,-m}$, and the echo is measured at the instant $2t_0$. When the HF signal is applied at the instant t_0, the coefficient α_{lm} reverses sign. The time variation of the spin density is then described by formulas (8), (10), and (11). Inasmuch as at the instant $2t_0$ the function $m_{lm}(\mathbf{c})$ returns to its initial value, we have at the echo time $2t_0$

$$m_{lm}(\mathbf{r}, 2t_0) = \int d\mathbf{q}\, m_{lm}(\mathbf{q}) \exp\left\{i\mathbf{q}\mathbf{r} - 2t_0\,\mathrm{Im}\left[\omega_{lm}^+(\mathbf{q} + \alpha_{lm}t_0\mathbf{k}) - \alpha_{lm}\int_0^{t_0} t\mathbf{k}\frac{\partial}{\partial \mathbf{q}}\omega_{lm}^+(\mathbf{q} + \alpha_{lm}t\mathbf{k})\,dt\right]\right\}. \tag{12}$$

In the long-wave limit $(m_{lm}(\mathbf{q}) \sim \delta(\mathbf{q}))$ this yields a decrease of the amplitude of the initial mode of the spin density by a factor

$$\ln\left|\frac{m_{lm}(\mathbf{r}, 0)}{m_{lm}(\mathbf{r}, 2t_0)}\right| = 2t_0\nu_l + \frac{2}{3}k^2(1 + \beta_l)(\Omega_0 + m\Omega)^2 t_0^3\,\mathrm{Im}\,D_{lm}(\mathbf{k}). \tag{13}$$

In deriving this formula, the function ω^+ was expanded in powers of k and formulas (4) and (6) were employed. The coefficient $D_{lm}(\mathbf{k})$ is determined by expression (6), in which it is necessary to replace the angles θ_1 and φ_1 by the angles θ_0 and φ_0 of the vector k. The first term in the right side of (13) represents the usual dissipative damping of the spin wave, while the second determines the effect of the gradient of the magnetic field. It is seen from (13) that it is best to observe the echo by exciting the mode $l = 0$, for in this case the first term turns out to be very small $(T \gg \tau_l)$, and the effect of the inhomogeneity of the magnetic field will become most clearly manifest. For the mode $l = 0$ we have

$$\mathrm{Im}\,D_{00}(\mathbf{k}) = \frac{v_F^2}{3}(\nu_1 - \nu_0)(1 + \beta_0)(1 + \beta_1)\left\{\frac{\cos^2\theta_0\,\bullet}{(\beta_0 - \beta_1)^2\Omega_0^2 + (\nu_1 - \nu_0)^2} + \frac{\sin^2\theta_0[(\beta_0 - \beta_1)^2\Omega_0^2 + \Omega^2(1 + \beta_1)^2 + (\nu_0 - \nu_1)^2]}{[(\beta_0 - \beta_1)^2\Omega_0^2 - \Omega^2(1 + \beta_1)^2]^2 + (\nu_0 - \nu_1)^2[(\nu_0 - \nu_1)^2 + 2\Omega_0^2(\beta_0 - \beta_1)^2]}\right. \tag{14}$$

The mode $l = 0$ is preferable from the point of view of experiment, since it is precisely this mode which determines the macroscopic magnetization, which is easiest to measure. For the higher modes $(l > 0)$, the dissipative damping of the spin waves will, generally speaking, prevail over the effect of the small gradient of the magnetic field, even if we disregard the difficulty of measuring higher moments of the spin density. It is seen from (13) and (14) that the damping of the mode $l = 0$ in a weakly-inhomogeneous magnetic field changes strongly when the temperature of the Fermi liquid is raised. Since $\nu_1 \sim T_M^2$ in very pure metals at low temperatures, there exists a temperature at which the imaginary part of the diffusion coefficient is maximal and the effect of the weakly-inhomogeneous magnetic field is most strongly manifest. An experimental measurement of the imaginary part of the diffusion coefficient can yield direct information concerning the interaction parameters β_l of a charged Fermi liquid.

We note that the approximation of homogeneous initial magnetization and the neglect of the finite dimensions of the sample are valid when the diffusion path of the electron during the observation time is small compared with the dimensions of the sample or with the dimension of the inhomogeneity of the HF signal. In the case of the ordinary skin effect with an HF signal whose wave vector is perpendicular to the constant magnetic field, the latter is satisfied when $\Omega_l > v_F/\delta t_0^{1/2}$, where δ is the depth of the skin layer. In pure metals, the mean free path of the electron along the magnetic field is $l_{\parallel} \sim 1$–3 mm, and if the observation time is $t_0 \sim 10^{-6}$ sec, then the required magnetic fields are $B \gtrsim 200$ kOe. Such fields have already been attained at present.

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