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Macroscopic tunneling of the oscillatory part of the magnetization at quantizing magnetic fields

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ABSTRACT

Tunneling of the oscillatory part of the magnetization is considered in small metallic samples at very low temperatures in quantized magnetic fields. By using the instanton approach, we calculate the action integral for the tunneling rate of the magnetization under conditions of the strong magnetic interaction between conduction electrons and the de Haas-van Alphen effect in specimens, which may undergo diamagnetic phase transitions. The obtained low tunneling temperature independent frequency may be detected by a resonance way. The temperature of the crossover between the classical and tunneling regimes of magnetization fluctuations is calculated.

1. Introduction

The existence of magnetic domains of the electron spin origin is well known in literature [1]. The phenomenon of domain formation associated with orbital magnetic moments has been detected in silver, beryllium, white tin, aluminium and lead [2–7]. These non-spin domains have been observed under conditions of the strong, nonlinear de Haas-van Alphen (dHvA) effect. One of the reasons for this behavior of metal samples is the instability of an electron gas called “diamagnetic phase transition” and leading to formation of Condon non-spin domains [3]. It is caused by magnetic interactions between conduction electrons, which become important when the internal field in the sample is significantly different from the applied magnetic field (the Shoenberg effect) [2]. As was shown by Condon, the above instability of the electron gas makes possible stratification of the sample into domains [8]. The state of art in this field has been reviewed in papers [4,9].

In recent years there has been considerable interest in the phenomenon of macroscopic quantum tunneling [10]. It corresponds to the tunneling of a macroscopic variable through the barrier between two minima of the effective potential of a system. The possibility of tunneling of the magnetic moment in nanometer-scale particles at very low temperatures has been detected and studied in ferromagnets and antiferromagnets [11–13].

The study of diamagnetic phase transitions observed in quantizing magnetic fields and at very low temperatures allows examining an additional quantum effect. It is of importance to investigate the above mentioned magnetic ordering in small, one domain, metallic samples exhibiting dHvA oscillations. The manifestation of the macroscopic tunneling in ferromagnets is the transition between two magnetization states of opposite directions. The effect happens, if the specimen consists of a single domain. In classical physics at high temperatures there are fluctuations resulting from thermal agitation. In quantum physics at low temperatures, the magnetic particle can move from one magnetic orientation to another by tunneling. Hence, the problem resembles quantum nucleation, i.e. the appearance of a nucleus of a domain of the orientation, opposite to the initial one. Tunneling of the oscillatory part of the magnetization, i.e., quantum fluctuations of the magnetization in metals may occur in the same way.

In this paper we compute the rate of tunneling of the orbital magnetization in small metallic samples. By using the instanton approach, we calculate the exponential of the action integral for the tunneling splitting induced by quantum dynamics of the oscillatory part of the magnetization.

In all the cases of the existence of this instability of the electron gas the dissipation effect on the tunneling rate seems to be negligible since the tunneling occurs at very low temperatures and in very pure samples, otherwise the diamagnetic phase transition is unable to occur.

2. Model

The oscillatory part of the density of the thermodynamic potential can be written neglecting all harmonics in the Lifshitz-Kosevich within the framework of the first harmonic approximation [2]:

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$$\Omega = \frac{1}{4\pi k^2}\left[a\cos(b) + \frac{1}{2}a^2\sin^2(b)\right], \tag{1}$$

where $b = k(B - H_0) = k[h_{ex} + 4\pi(1 - n)M]$, M is the oscillatory part of the magnetization, H_0 is the magnetic field inside the material, ($k = \frac{2\pi F}{H_0^2}$, F is the fundamental frequency of dHvA oscillations, n is the demagnetization factor,$h_{ex} = H_{ex} - H_0$, is the small increment of the external magnetic field, B is the magnetic induction).

In the first harmonic approximation, the magnetization is found from the implicit equation of state [2]:

$$4\pi kM = a\sin[k(h_{ex} + 4\pi(1 - n)M)], \tag{2}$$

where $a = 4\pi(\frac{\partial M}{\partial B})_{B=H_0}$ is the reduced amplitude of magnetization oscillations. If $a > 1$, a state of lower density of the thermodynamic potential can be achieved over part of an oscillation cycle by the sample breaking up into domains. Since the demagnetization factor n is not a well-defined quantity we omit it in the further use of Eq. (2). However, the demagnetization problem will be further considered. We examine the behavior of the magnetization in the center of the period of dHvA oscillations, where the increment of the external magnetic field is equal to zero: $h_{ex} = 0$. This case corresponds to ferroelectrics and ferromagnets at zero external field leading to the appearance of 180^0 domains.

Close to the phase transition temperature (which is found from the equation $a = 1$ in the case of infinite specimens) we can present the density of the thermodynamic potential as an expansion in powers of the magnetization in the center of dHvA oscillations [9]:

$$\Omega = 2\pi(1 - a)M^2 + \frac{8}{3}\pi^3k^2aM^4. \tag{3}$$

Therefore, we arrived at the Landau-type thermodynamic potential:

$$\Omega = -\frac{\alpha}{2}M^2 + \frac{\beta}{4}M^4, \tag{4}$$

$\alpha = 4\pi(a - 1)$, $\beta = \frac{32}{3}a\pi^3k^2$. According to [2], the temperature and magnetic field dependence of the reduced amplitude of dHvA oscillations is:

$$a(T, H, T_D) = a_0(H)\frac{\lambda T}{\sinh(\lambda T)}\exp(-\lambda T_D), \tag{5}$$

$\lambda \equiv \frac{2\pi^2k_Bm_cc}{ehH}$, m_c is the cyclotron mass, $\eta_c = \frac{m_c}{m}$, m is the electronic mass, c is the velocity of light, k_B is the Boltzmann constant, e is the absolute value of the electron charge, $\hbar = \frac{h}{2\pi}$, h is the Planck constant, T_D is the Dingle temperature, $a_0(H) = (\frac{H_{\max}}{H})^{\frac{3}{2}}$ is determined by $H_{\max}$ which is the largest field at which the domain phase exists. Taking into account the inhomogeneous term $\frac{1}{2}K(\frac{\partial M}{\partial x})^2$ [14] (K is the positive coefficient of the inhomogeneous term, proportional to r_c^2, where r_c is the cyclotron radius) and including the demagnetization energy per unit volume $\frac{14}{\pi^2}\zeta(3)M_0^2D/L$ [15], where ζ is the ζ- function, M_0 is the magnetization in the center of the domain, D is the domain width, minimizing the density of the thermodynamic potential with respect to the magnetization in the domain center and domain width $\frac{\partial\Omega}{\partial M_0} = 0$ and $\frac{\partial\Omega}{\partial D} = 0$, and using the surface energy $\sigma = \frac{4}{3}\frac{KM_0^2}{d}$ [16], where d is the domain wall thickness, $d = \frac{r_c}{(a-1)^{\frac{1}{2}}}$ [9] and $K = \frac{r_c^2}{4}$, we derive a shift of the diamagnetic phase transition compared to the infinite sample case and a shift of the temperature-magnetic field phase diagram [17].

Equation $a = 1$ describes the temperature-magnetic field phase diagram determining the locus of points showing the series of the phase-transition temperatures in an infinite sample. The envelope of these points is the diamagnetic phase-transition boundary. However, as it follows from [17], the domain phase appears in the slab not at $a = 1$ but at

$$a = 1 + (\frac{1}{3})^{\frac{4}{3}}\left(\frac{r_c}{L}\right)^{\frac{2}{3}}, \tag{6}$$

in the film of thickness L. Thus, the phase transition occurs at

$$a^* = 1, \quad a^* = a - \left(\frac{1}{3}\right)^{\frac{4}{3}}\left(\frac{r_c}{L}\right)^{\frac{2}{3}} \tag{7}$$

3. Results and their discussion

We consider the double quantum effect. Firstly, the magnetic field should be a quantized one for the existence of the dHvA effect. The second quantum effect is the macroscopic tunneling of the oscillatory part of the magnetization. The tunneling mode should be independent of temperature. Thus, the reduced amplitude of magnetization oscillations a^* does not depend on temperature. This occurs at very low temperatures for which $\lambda T << 1$ (Eq. (5)).

In the center of the period of dHvA oscillations the density of the thermodynamic potential (4) has two degenerated energy minima. The tunneling removes the degeneracy of the ground state, corresponding to two opposite directions of the magnetization, leading to the new ground state, which is a superposition of the up and down states. In the semi-classical limit, the tunnel splitting Δ can be calculated using standard methods like WKB [18] or the instanton technique [19] to obtain analytical results. In the latter, the problem of finding Δ is essentially reduced to determine the classical path, which connects the two maxima of the inverted potential $-\Omega(M)$ (the minima of $\Omega(M)$) and calculate the corresponding action. Actually, the tunneling of magnetization should be described as a sum over all possible paths of transitions from the initial state to the final one, including those that are classically forbidden. Among the paths, the saddle point path dominates. It passes through the valley of the energy, i.e., the thermodynamic potential density relief. For a quartic double-well potential of the form (4) these calculations can be performed analytically. Using [19], we obtain

$$\Delta = 4\sqrt{3}\omega_0\left(\frac{E}{2\pi}\right)^{\frac{1}{2}}\exp(-E), \tag{8}$$

$$E = \frac{S_I}{\hbar}, \tag{9}$$

$$S_I = \frac{16V\alpha^2}{3\omega_0\beta}, \tag{10}$$

where S_I is the instanton action obtained from the sum over paths, V is the volume involved in the tunneling process given by $V = L_xL_yL_z$, the applied magnetic field is oriented along the axis z, the sizes of L_x and L_y are taken to be close to the thickness of a slab sample $L_z = L$, $L_x = L_y \approx L$. The pre-factor ω_0 is a classical attempt frequency related to the cyclotron frequency ω_c by the formula $\omega_0 = \omega_c(a^* - 1)^{1/2}$. The presented mode softening is determined by the magnetic interactions between conduction electrons leading to Condon domains. The factor $(a^* - 1)^{1/2}$ is the result of the scaling procedure characteristic of the mean field theory.

The tunneling splitting value,Δ, is given by Eqs. (7)–(10). Taking into account Eqs. (7)–(10) we obtain:

$$E = \frac{8(a^* - 1)^{3/2}\gamma^2V}{a^*\hbar\omega_0}, \tag{11}$$

$$\gamma^2 = \frac{H^4}{4\pi^2F^2}, \tag{12}$$

γ is the period of the dHvA oscillations.

We estimate the tunneling splitting (8). For instance, for silver, taking $T_D = 0.17K$,$H = 37T$, $L = 240nm$, $a^* = 1.01$ we obtain the tunneling splitting $\Delta = 1.9{\cdot}10^3Hz$ close to those observed in magnetic metals