Quantum tunneling of magnetization in a small-area domain wall

B. A. Ivanov and A. K. Kolezhuk
Institute of Metal Physics, Ukrainian National Academy of Sciences, 252142 Kiev, Ukraine

(Submitted 4 November 1994)
Pis'ma Zh. Eksp. Teor. Fiz. **60**, No. 11, 792–795 (10 December 1994)

A domain wall of small area in an antiferromagnet is a good setting for observing a macroscopic quantum tunneling of magnetization.

Macroscopic quantum tunneling in magnetic materials has recently attracted considerable interest, both theoretical and experimental (see a review[1]). Macroscopic quantum tunneling in magnetic materials corresponds to below-barrier transport between macroscopic equilibrium states of the magnetization distribution. This phenomenon has been studied theoretically for mesoscopic systems—ferromagnetic[2] and antiferromagnetic[3] particles of ultrasmall size (~50–100 Å)—and also for the depinning of domain walls.[4] Effects of macroscopic quantum tunneling are manifested in a finite magnetic relaxation as $T \to 0$ (see the bibliography in Ref. 1) and also in the onset of a resonance involving levels split by tunneling in ultrasmall particles[5] (see also the discussion in Ref. 6).

1. In this letter we suggest a new effect of macroscopic quantum tunneling: a below-barrier quantum change in the polarity of a small-area domain wall. We consider a wall in a thin magnetic film which is a narrow strip or "hourglass," with the wall localized at the neck. (Yet another possibility arises from the use of 2D magnetic materials such as films of manganese stearate.[7]) We assume that the magnetic material is orthorhombic with an anisotropy energy

$$w_a = K_1 n_y^2 + K_2 n_z^2, \quad K_2 > K_1, \tag{1}$$

where $\mathbf{n}$ is a unit vector (the normalized magnetization in the case of a ferromagnet; the antiferromagnetism vector in the case of an antiferromagnet), K_1 and K_2 are anisotropy constants, XY is the easy plane, and Y is the easy axis in that plane. In terms of the angle variables $n_y = \cos\theta$, $n_x + i n_z = \sin\theta \exp i\varphi$, a wall corresponds to the solution

$$\cos\theta = \tanh(x/\Delta), \quad \varphi = \varphi_0 = \text{const}, \tag{2}$$

where $\Delta = \Delta_0 \equiv (A/K_1)^{1/2}$, A is the inhomogeneous-exchange constant, and we have $\varphi_0 = 0, \pi$ for a suitable wall [here we are using (1)]. The wall can thus be in two states, which differ in the direction of $\mathbf{n}$ at the center of the wall (at $x=0$) : $\mathbf{n}(0) = +\mathbf{e}_x$ and $\mathbf{n}(0) = -\mathbf{e}_x$ (Fig. 1). The energies of these states are the same, $S_0\sigma$, where S_0 is the area of the wall, and $\sigma = 4(AK_1)^{1/2}$ is the energy of the wall per unit area.[8] These states are separated by a finite barrier $U_0 = 4S_0\Delta_0[(K_1K_2)^{1/2} - K_1]$, so a tunneling can occur between them. In either a ferromagnet or an antiferromagnet with a Dzyaloshinskiĭ interaction, polarization reversal of the wall is accompanied by a change in the projection of the

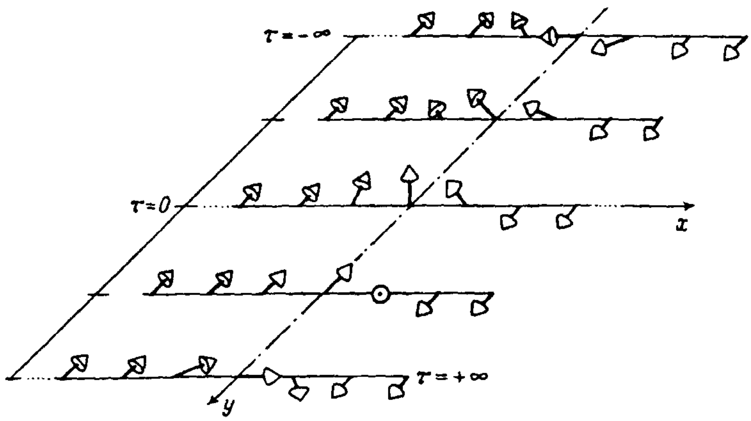

FIG. 1. Distribution (shown schematically) of bfn in a domain wall in the course of polarization reversal. The labels $\tau=-\infty$ and $\tau=+\infty$ correspond to equilibrium states: an initial state with $\mathbf{n}(0)=-\mathbf{e}_x$ and a final state with $\mathbf{n}(0)=+\mathbf{e}_x$. The value $\tau=0$ corresponds to the state $\mathbf{n}(0)=+\mathbf{e}_z$, with the maximum energy. It is easy to see that, in the x,τ plane, the process corresponds to a vortex configuration in which the vector $\mathbf{n}$ rotates through 2π as a closed loop around the point $x=0,\ \tau=0$ is traversed.

magnetic moment of the wall onto an axis [the X axis for a ferromagnet; the ($\mathbf{e}_x\times\mathbf{d}$) axis for an antiferromagnet (see the discussion below)] by an amount $\Delta M=MS_0\Delta_0$, where M is the magnetization per unit volume. From the standpoint of the "classical" parameters ΔM and U_0, the wall thus behaves as a small particle with a volume $S_0\Delta_0$.

We have analyzed the tunneling probability in a wall for the cases of ferromagnets and antiferromagnets. For ferromagnets, this probability turns out to be the same as for a small particle of volume $S_0\Delta_0$, i.e., for a particle which has one linear dimension proportional to the macroscopic quantity $\Delta_0\sim(10^2-10^3)a$, where a is the constant of the magnetic lattice. However, the case of an antiferromagnet, discussed below, is of considerably more interest. We will show that in this case the probability for polarization reversal does not depend on the wall thickness, and that it is determined exclusively by the wall area S_0.

2. We describe the dynamics of the antiferromagnet on the basis of an effective Lagrangian for a unit antiferromagnetism vector[8] $\mathbf{l}$:

$$L=\int d^3x\{A[(1/c^2)(\partial\mathbf{l}/\partial t)^2-(\nabla\times\mathbf{l})^2]-K_1l_y^2-K_2l_z^2\}, \tag{3}$$

where c is the limiting velocity of spin waves. The magnetization of the antiferromagnet, $\mathbf{M}$, is determined by the Dzyaloshinski ĭ field $\mathbf{d}$: $\mathbf{M}=\chi_\perp(\mathbf{d}\times\mathbf{l})$, where $\chi_\perp$ is the susceptibility of the antiferromagnet. On the basis of (3) we can describe a "rocking" mode localized at a wall. In the linear approximation, that mode corresponds to

$$\varphi=\varphi_0+\psi_0\cos\omega_l t,\quad \omega_l=\omega_0\sqrt{\rho},\quad \rho=K_2/K_1-1, \tag{4}$$

where the quantity $\omega_0 = c/\Delta_0$ is the same as the activation energy for the lower branch of spin waves of the antiferromagnet. Excitation of this mode has been observed in orthoferrites.[9]

It is fairly clear that this localized mode is responsible for the tunneling. Assuming that the area of the wall is small, we restrict the discussion to an analysis of 1D field configurations $\mathbf{l} = \mathbf{l}(x,t)$. To analyze rocking with an amplitude which is not small, we seek a solution in the form of (2) with $\varphi = \varphi_w(t)$, $\Delta = \Delta_0(1 + \rho \sin^2 \varphi_w)^{-1/2}$ where we can treat φ_w as a slow variable under the condition $\omega_l \ll \omega_0$ ($\rho \ll 1$). Using (3), we then find the effective Lagrangian for the one variable φ_w in the leading approximation in ρ, $\omega_0^{-1} d\varphi_w/dt$:

$$L = (1/2)\sigma S_0 \{\omega_0^{-2}(d\varphi_w/dt)^2 - \rho \sin^2 \varphi_w\}. \qquad (5)$$

This Lagrangian describes a "particle on a circle" in a two-well potential. The magnitude of the tunneling splitting can be calculated easily either by instanton methods[10] or by the WKB approximation. The tunneling matrix element is $\hbar\Gamma \simeq (\hbar\omega_c/2\pi)\exp(-S_E/\hbar)$, where ω_c is a characteristic frequency of the classical motion, and S_E is the Euclidean action calculated for the classical solution in the imaginary time $\tau = it$, which couples the initial and final states. In our case, solutions of this type are instantons and antiinstantons of the type $\varphi_w(\tau) = \pm 2\arctan\exp(\omega_0\sqrt{\rho}\tau)$, and the magnitude of the tunneling splitting of the levels, $\Delta E = 2\hbar\Gamma$, is given by

$$\Delta E = \frac{2}{\pi}\,\omega_0\sqrt{\rho}\exp\left\{-2S_0\frac{\sigma}{\hbar\omega_0}\sqrt{\rho}\right\}. \qquad (6)$$

It is a simple matter to calculate the critical temperature T_c, below which effects of macroscopic quantum tunneling outweigh thermal effects, by writing the exponential function in (6) as $\exp(-U_0/kT_c)$, where U_0 is the barrier height. As a result, we find the estimate $kT_c \sim \hbar\omega_0\sqrt{\rho}$. Dissipation effects may cause a renormalization[11] of T_c. In the case of a higher symmetry (e.g., tetragonal), we would have obtained a potential with more minima, but the picture would remain fundamentally the same.

3. Let us compare (6) with the corresponding result for macroscopic quantum tunneling in a small antiferromagnetic particle (for the latter case the argument of the exponential function is $KV_0/\hbar\omega_0$, where V_0 is the volume of the particle[3]). Interestingly, the argument of the exponential function in (6) is $\sqrt{\rho}S_0 s/a^2$ in order of magnitude, where s is the spin of the magnetic atom; it is independent of both the exchange integral J and the wall thickness Δ_0. In other words, it is proportional to the number of magnetic atoms in the cross section of the wall, not the number in the volume of the wall. The observable parameter ΔM has the behavior $\Delta M \propto \Delta_0$, and increases with increasing wall thickness. The temperature $T_c \propto \hbar\omega_0$ increases as $\sqrt{J}$ with increasing value of the exchange integral. Since ΔM is not small, there is the hope that macroscopic quantum tunneling of a domain wall in an antiferromagnet might be detected not only as a resonance involving split levels (dissipative effects might be a problem here[1]) but also through direct measurement of the magnetization with a SQUID magnetometer.

Furthermore, in our case the argument of the exponential function is independent of the uniaxial anisotropy constant ($\sigma \propto \sqrt{K}$ and $\hbar\omega_0\sqrt{K}$). It is proportional to $\sqrt{\rho}$, where

ρ determines the ratio of the anisotropy in the basal plane to the uniaxial anisotropy. For easy-axis antiferromagnets with a principal axis C_n, $n=3,4,6$, the quantity ρ is small. For tetragonal MnF_2, for example, the fields of the uniaxial and intraplane anisotropy are 30 Oe and 7.8 kOe, respectively,[12] and we have $\rho=3.8\times10^{-3}$. For orthorhombic antiferromagnets of the orthoferrite type, we have values $\rho\simeq0.1-1$, but these values may decrease substantially if a magnetic field is applied along an intermediate axis. In this case we would have $\rho\to0$ as $H\to H_c$, $\chi_\perp H_c^2=K_2-K_1$. In the limit $\rho\to0$, however, there is a decrease in the pre-exponential factor and in the value of T_c.

4. Let us carry out an estimate for a typical antiferromagnet. For definiteness we select MnF_2, with $\chi_\perp\simeq1.6\times10^{-4}$, $\sigma\simeq0.5$ erg/cm^2, $H_{SF}\simeq93$ kOe, and $\omega_0\simeq1.6\times10^{12}$ s^{-1}, i.e. $\hbar\omega_0/k\simeq12$ K. For these values, the argument of the exponential function is written in the form $S_0\sqrt{\rho}(6\times10^{14}\ \mathrm{cm}^{-2})$. Assuming that a macroscopic quantum tunneling is plausible if the argument of the exponential function does not exceed[1] 20–30, even with $\rho\simeq0.1$ we find $\sqrt{S_0}\leqslant4\times10^{-7}$ cm and $T_c\simeq4$ K. These figures are no worse than for macroscopic quantum tunneling in small antiferromagnetic particles. If we instead use $\rho\simeq4\times10^{-3}$, we find $\sqrt{S_0}\leqslant10^{-6}$ cm with $T_c\simeq0.8$ K. These results look realistic for an experimental observation of the effect.

We are indebted to E. M. Chudnovskiĭ for stimulating discussions. This study was supported in part by Grant UB7000 from the International Science Foundation and Grant 2.2/36 from the Ukrainian State Committee on Science and Technologies.

[1] E. M. Chudnovsky, J. Appl. Phys. **73**, 6697 (1993).
[2] E. M. Chudnovsky and L. Gunther, Phys. Rev. Lett. **60**, 661 (1988); Phys. Rev. B **37**, 9455 (1988).
[3] B. Barbara and E. M. Chudnovsky, Phys. Lett. A **145**, 205 (1990); I. V. Krive and O. B. Zaslavskii, J. Phys. Cond. Matt. **2**, 9457 (1990).
[4] P. C. E. Stamp, Phys. Rev. Lett. **66**, 2802 (1991); E. M. Chudnovsky *et al.*, Phys. Rev. B **46**, 5392 (1992).
[5] D. D. Awschalom *et al.*, Phys. Rev. Lett. **65**, 783 (1990), **68**, 3092 (1992).
[6] Anupam Garg, Phys. Rev. Lett. **70**, 2198 (1993), **71**, 4249 (1993); D. D. Awschalom *et al.*, Phys. Rev. Lett . **70**, 2199 (1993); **71**, 4276 (1993).
[7] M. Pomerantz, Surf. Sci. **142**, 556 (1984).
[8] V. G. Bar'yakhtar *et al.*, Usp. Fiz. Nauk **146**, 417 (1985) [Sov. Phys. Usp. **28**, 563 (1985)].
[9] G. A. Kraftmakher *et al.*, Zh. Eksp. Teor. Fiz. **63**, 1353 (1972) [Sov. Phys. JETP **36**, 714 (1972)].
[10] R. Rajaraman, *Solitions and Instantons: An Introduction to Solitons and Instantons in Quantum Field Theory* (Elsevier, 1982).
[11] A. J. Leggett, Rev. Mod. Phys. **59**, 1 (1987).
[12] Y. Shapiro and T. Zak, Phys. Rev. **170**, 503 (1968).

Translated by D. Parsons