Magnetoacoustic resonance of domain boundaries in antiferromagnets

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The magnetoelastic.oscillations of the domain boundary in an easy-axis antiferromagnet; which are attributable to nonuniform, magnetostrictive deformations in it, were examined. New resonances of the domain boundary, including a low-frequency resonance characteristic of an easy-plane antiferromagnet, are predicted.

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The so-called magnetoelastic gap (ME gap) effects[1,2] (or "solidified lattice" effects[3]), which are caused by the influence of spontaneous, magnetostrictive deformations in the ground state of a magnetically ordered system on the spectrum of coupled magnetoelastic oscillations, exist in homogeneous (single-domain) ferro- and antiferromagnets.

For both the physics of magnetism as well as physics generally it is of interest to consider the analogous effects for the ME oscillations of domain boundaries. The latter destroy not only the rotational (as in the single-domain case) but also the translational symmetry. The point of this discussion is to consider the influence of inhomogeneous (along the normal to the domain boundary), spontaneous, magnetoelastic deformations on the coupled boundary oscillation and sound spectrum.

Assuming the existence of a domain boundary,[4] in which the vector $\mathbf{L}=\mathbf{M}_1-\mathbf{M}_2$ is turned from the direction $\mathbf{L}\uparrow\uparrow Z$ (polar angle $\theta=0$) for $y=-\infty$ to the direction $\mathbf{L}\uparrow\downarrow Z$ ($\theta=\pi$) for $y=+\infty$ and also assuming that that azimuth angle ϕ remains constant, from the condition of minimum total energy of the easy-axis antiferromagnet we first find the equilibrium distribution of the local sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$ (determined by the moduli $|\mathbf{M}_1|=|\mathbf{M}_2|=\mathbf{M}_0$ and by the angles θ_i, ϕ_i) and of the deformations u_{ik} corresponding to a Bloch boundary (rotation in the XZ plane) and a Néel boundary (rotation in the YZ plane). In particular, for the latter[1)]:

$$\theta_1=\pi-\theta_2=\theta,\quad \phi_1=\phi_2-\pi=\pi/2,\quad u_{xy}=u_{xz}=0,$$
$$u_{xx}=u_{xx}(\pm\infty)=u_{yy}(\pm\infty)\equiv U_\perp,\quad u_{zz}=u_{zz}(\pm\infty)\equiv U_{\parallel},\tag{1}$$

$$u_{yz}=-\frac{B_{44}}{4C_{44}}\sin 2\theta,\quad u_{yy}=U_\perp-\frac{B_{\parallel}}{C_{\parallel}}\sin^2\theta,\tag{2}$$

$$\operatorname{ctg}\theta=-\sqrt{d}\,\operatorname{sh}\frac{y}{\Delta},\quad d=1+\frac{K_2^*}{K_1^*},\quad \Delta^2=\frac{A}{K_1^*},\tag{3}$$

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where $C_{\alpha\beta}$ and $B_{\alpha\beta}$, the elastic and magnetoelastic constants, are labeled by two indices in the standard manner; A is the exchange constant for $L^{-2}(\partial L/\partial x_k)^2$, K_1^* and K_2^* are the renormalized magnetostrictive constants of the anisotropy (for L_z^2 and L_z^4). Let us concentrate on the existence of the inhomogeneous deformations (2) at the domain boundary.

If the linearized equations of motion are now written in the usual manner for the homogeneous, coupled oscillations of the magnetizations ($\Delta M_{1,2}$) and displacements (Δu_i) in the XZ plane near the ground state (1)–(3), then it can be shown that they break up into two independent subsystems, respectively, for the variables:

$$\text{I.}\quad \Delta M_x,\quad \Delta L_y,\quad \Delta L_z,\ \Delta u_y,\ \Delta u_z, \tag{4}$$

$$\text{II.}\quad \Delta L_x,\quad \Delta M_y,\quad \Delta M_z,\quad \Delta u_x\ (\mathbf{M} = \mathbf{M}_1 + \mathbf{M}_2)\,. \tag{5}$$

The variables from set II become zero for the type I oscillation mode, and vice versa. The approximate solution of these subsystems leads to dispersion equations of the form:

$$\text{I.}\quad \omega^2[\Phi(\omega) - 1] = 1, \tag{6}$$

$$\text{II.}\quad \omega \approx \omega_{ME} - i\Gamma, \tag{7}$$

where

$$\Phi(\omega) = \beta_1 \int \frac{x^3}{\mathrm{sh}^2 x} \frac{dx}{\dfrac{\omega}{\omega_y} - x + i0} + \beta_2 \int \frac{x^3}{\mathrm{ch}^2 x} \frac{dx}{\dfrac{\omega}{\omega_z} - x + i0} \equiv \Phi_1(\omega) - i\Phi_2(\omega), \tag{8}$$

$$\beta_1 = \frac{\omega_{ME1}^2}{\omega_y^2};\qquad \beta_2 = \frac{B_{44}}{8B_{11}}\frac{C_{11}}{C_{44}}\beta_1;\qquad \omega_{ME1}^2 = \frac{\gamma^2 A_0 B_{11}^2}{2C_{11}M_0^2} \tag{9}$$

(A_0 are the homogeneous exchange constants for M^2),

$$\omega_{ME} \approx \frac{\gamma^2 A_0}{3M_0^2}\left[\frac{4B_{44}}{C_{44}} + \frac{4B_{11}(B_{11} - B_{12})}{C_{11}} - \frac{(B_{11} - B_{12})^2}{C_{11} - C_{12}}\right], \tag{10}$$

$$\frac{\Gamma}{\omega_{ME}} \approx \frac{\omega_{ME2}^2\,\omega}{\omega_{ME}^2\,\omega_x};\qquad \omega_{ME2}^2 = \frac{\gamma^2 A_0 (B_{11} - B_{12})^2}{4(C_{11} - C_{12})M_0^2}, \tag{11}$$

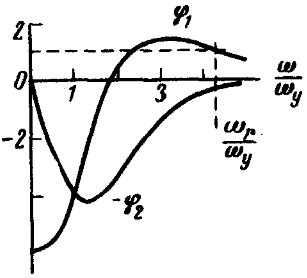

FIG. 1.

$$\omega_i = s_i/\Delta, \quad s_x = \sqrt{(C_{11} - C_{12})/\rho}, \quad s_y = \sqrt{C_{11}/\rho}, \quad s_z = \sqrt{2C_{44}/\rho}. \tag{12}$$

Equations (8)–(12) are given for the simplest special case $d = 1$ ($K_2^* = 0$). The more general case, which will be discussed in another place, constains no fundamental differences (if one excludes the extreme limit of very small d values).

The $\omega = 0$ root of Eq. (6) corresponds to the Goldstone mode which is associated with degeneracy of the ground state relative to equilibrium translations of the domain boundary. The roots of the equation

$$\Phi_1(\omega) = 1, \tag{13}$$

which follows from (6) after the splitting off of the Goldstone mode, give quasilocal (or resonance) modes corresponding to relative oscillations of the domain boundary and the inhomogeneous deformations created by them. As an illustration, Fig. 1 depicts schematically the graphical solution of Eq. (13) for the case $\beta_2 \ll \beta_1$. Real roots exist only for the condition $\Phi_1(\omega_{\max}) > 1$ or $\beta_1 > 1$ in order of magnitude, which has a simple physical meaning according to (9): the tuning frequency ω_y of the deformations per thickness Δ is less than the characteristic magnetoelastic frequency ω_{ME_1}. Unfortunately, this can occur only in antiferromagnets than have certain extreme parameters (for example, for $B_{\alpha\beta} \sim 10^8$ ergs/cm^3, if $\Delta \sim 10^{-6}$ cm, $A_0/M_0 \sim 10^6$ Oe, $M_0 \sim 10^2$ Oe). These oscillations are damped, since upon their excitation (by a varying magnetic field directed along the normal to the boundary–the Y axis), according to (4), elastic waves must be generated with a wavelength of the order of Δ. The damping is described by the imaginary part of Eq. (8), which is also shown in Fig. 1. It can be seen in Fig. 1 that only the upper mode (ω_r), for which the damping can be quite small, is a resonance mode.

The other resonance mode II (7), whose frequency is determined by Eq. (10) and the damping by Eq. (11), is obviously of greater interest. This mode is excited by a varying field parallel to the Z axis, and in this case a transverse sound Δu_x, which diverges along the normal from the boundary, is generated. If, on the other hand, this sound is incident on the boundary, then at the frequency (10) it will be scattered by it in a resonance manner. We note that the frequency (10) is the MOE gap for the corresponding branch of the "intraboundary" (two-dimensional) spin waves, which is analogous to that for volume spin waves.[2]

The experimental observation of the predicted magnetoelastic resonances of the domain boundary would definitely be of interest as one specific example of a breakdown of translational symmetry in a system of coupled waves of a different type.

[1)]From the viewpoint of the effects of interest to us, similar results are obtained for a Bloch boundary, which we do not examine here.

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