MAGNETOELASTIC COUPLING AS A SOURCE OF SHAPE DEPENDENCE OF AFMR SPECTRA

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UDC 532

We study the possible influence of the crystal shape on AFMR spectra in the framework of a phenomenological theory with regard for the destressing energy. It is shown that, for the crystals with strong magnetoelastic coupling, the crystal shape may be a source of artificial anisotropy of the magnetoelastic nature. The shape-induced anisotropy may be greater than the bare magnetic anisotropy of a crystal. If this is the case, the gap in AFMR spectra must be sensitive to the orientation of an external magnetic field.

1. Introduction

Antiferromagnetism as a phenomenon was discovered nearly a half-century ago by L. Néel, L.D. Landau, L.V. Shubnikov [1, 2] but it was treated for a long time only from the fundamental (unpractial) point of view. The reason for such an attitude to antiferromagnetic materials (AFM) stems from the absence of a macroscopic magnetization that complicates the application of AFM in, e.g., information-carrying media or control devices. Nevertheless, the magnetic, magnetoelastic, static, and dynamic properties of antiferromagnets were thoroughly studied by the disciples of Landau's phenomenological methods.

An important contribution to the development of the theory of antiferromagnets was made by researches of the Ukranian branch of Landau's school headed by A.I. Akhiezer and personally by one of the prominent representatives of this school, Academician V.G. Bar'yakhtar. In connection with the 75th anniversary of V.G. Bar'yakhtar, it is of our pleasure to mention just his achievements and the achievements of his numerous disciples in the field of antiferromagnetic phenomena: phenomenological approaches to the description of orientational phase transitions [3, 4] and the dynamics of magnetic perturbations, expressions for AFM wave spectra and coupled magnetoelastic oscillations in surfaceless crystals and thin plates, calculations of domain wall profiles in AFM, etc. Together with a lot of interesting and important results made by other researches, they laid the foundations and principles of the phenomenological theory of antiferromagnetism at the end of the 1980s.

Nevertheless, at the end of the 1990s, the interest to AFM crystals suddenly increased due to new achievements in the technology and application of thin films and multilayers for the information storage. It was found that AFM materials, though being inert from the magnetic point of view, can be very promising systems in combination with ferromagnetic (FM) materials for high-density information media. Another attractive feature of AFM is very high (as compared to FM) magnetostriction, which makes these materials useful for the controller and actuators driven by an external magnetic field.

The technological application of AFM sets a new research problem, namely, to find out a mechanism of the observed macroscopic behaviour of AFM crystals in an external magnetic field and to develop a clear and convenient formalism for the description of these phenomena. Up to now, there is no common point of view concerning the origin and mechanism of the domain structure formation in compensated AFM. Experimental observations show that the domain structure really exists and may be easily and reversibly changed by application of an external field. On the other hand, because of the absence of a net magnetization, AFM crystals have no evident source of the long-range demagnetization field which is generally considered to be a cause of the domain structure formation in FM.

Nevertheless, the strong magnetoelastic coupling and the related phenomena such as coupled magnetoelastic oscillations [5], magnetoelastic phase transitions [4], intermediate state consisted of FM/AFM domains [6] made it possible to suggest that the domain structure in AFM has magnetoelastic origin. Such an idea was recently formulated by S.M. Ryabchenko et al. [7,8], and we [9, 10] further assumed the presence of the so-called

destressing energy in finite-size AFM samples analogous to the demagnetization energy in FM. The correctness of our hypothesis is validated by the agreement between the calculated and experimentally observed temperature and field dependences of magnetostriction [11] and magnetoresistance [12–15] of layered AFM and by the direct observation of AFM domains and domain wall profiles [16–18]. Nevertheless, these experimental and theoretical results are insufficient in order to finally and unambiguously confirm or disprove the models based on the concept of destressing fields.

In this paper, we suggest a method for a direct experimental proof of the adequacy of the developed model of domain structure formation. We propose to study experimentally the influence of a crystal shape on AFMR spectra for the compensated AFM with strong magnetoelastic coupling. We dedicate this paper to the 75th anniversary of Academician V. G. Bar'yakhtar, whose constribution to the development of the theory of magnetism and antiferromagnetism is widely acknowledged.

2. Internal Stresses. Destressing Energy

The phenomenological model developed in the present paper is based on the assumption that the AFM ordering is accompanied by appearance of the internal mechanical stresses. In fact, we suppose a certain analogy between AFM and martensitic (thermoelastic) phase transitions. Actually, in both cases, the primary order parameter (AFM vector $\mathbf{l}$ or, probably, some combinations of electronic functions) is a source of the so-called (according to [19]) quasiplastic internal stresses described by the second-rank tensor $\hat{\sigma}^{(\mathrm{in})}[\mathbf{l}(\mathbf{r})]$. Eigenvalues of this tensor (i.e. those characteristics which are invariant with respect to coordinate rotations) depend only on magnetoelastic constants $\hat{\lambda}$, while eigenvectors are rigidly coupled with the orientation of the AFM vector $\mathbf{l}$. The corresponding deformations should be determined from the minimum conditions for the crystal free energy with a due account of the magnetic anisotropy $f^{(\mathrm{mag})}$. On the other hand both factors, namely the bare magnetic anisotropy and the state of the elastic subsystem define, at the equal foot, the equilibrium orientation of the AFM vector as may be seen from the system of equations

$$\frac{\delta f^{(\mathrm{mag})}}{\delta l_j} = \frac{\partial \sigma^{(\mathrm{in})}_{jk}}{\partial l_m}\frac{\partial u_k}{\partial r_m}, \tag{1}$$

$$\frac{\partial}{\partial r_k} c_{jklm}\frac{\partial u_m}{\partial r_l} = \frac{\partial}{\partial r_k}\left(\sigma^{(\mathrm{in})}_{jk}[\mathbf{l}(\mathbf{r})]\right) \tag{2}$$

which is derived from the minimum conditions for the thermodynamic potential

$$\Phi[\mathbf{l}(\mathbf{r}), \hat{u}(\mathbf{r})] =$$

$$= \int\limits_V \left\{ f^{\mathrm{mag}}[\mathbf{l}(\mathbf{r})] + \frac{1}{2}\hat{u} : \hat{c} : \hat{u} - \hat{\sigma}^{(\mathrm{in})}[\mathbf{l}(\mathbf{r})] : \hat{u} \right\} d\mathbf{r}. \tag{3}$$

Here, $\hat{c}$ is the four-rank tensor of elastic moduli, $\hat{u}$ is the strain tensor, $\mathbf{u}$ is a displacement vector, and V is the crystal volume.

In the infinite (surfaceless) crystal, the transition-induced internal stresses [Eq. (2)] give rise to the well-known renormalization of magnetic anisotropy constants and the formation of the so-called isotropic magnetoelastic gap in AFM spectra [20, 21] but have no effect on the equilibrium orientation of the AFM vector.

In an opposite case of finite-size samples, Eqs. (1) and (2) should be solved with account of the (nontrivial) boundary conditions

$$n_k c_{jklm}\frac{\partial u_m}{\partial r_l}(\mathbf{r}_S) = \sigma^{(\mathrm{in})}_{jk} n_k(\mathbf{r}_S), \quad (\mathbf{r}_S \in S), \tag{4}$$

where $\mathbf{n}$ is the normal to the surface S at a point $\mathbf{r}_S$. According to our assumption, stress tensor components $\hat{\sigma}^{(\mathrm{in})}$ are tightly bound to the orientation of the AFM vector $\mathbf{l}$. So, $\hat{\sigma}^{(\mathrm{in})}$ may be nonzero at the sample surface and, hence, the surface orientation with respect to the crystal axes should have an influence on the equilibrium distribution of magnetic moments in the sample (see Eq. (4)). It should be stressed that, due to a long-range character of the elastic forces, this influence is quite strong. The corresponding free energy contribution is proportional to the crystal volume (see Eq. (5)) and, hence, may compete with the bare magnetic anisotropy.

Formally, the shape-induced effect may be described with the use of an additional term in the thermodynamic potential [10,22]. This term called the destressing energy $\Phi^{\mathrm{dest}}_{\mathrm{dd}}$ is analogous to the energy contribution of the demagnetization field which is responsible for the shape-induced anisotropy in ferromagnets. In the general case of a finite-size sample that has a volume V, the destressing energy may be represented in a simple form

$$\Phi^{\mathrm{dest}}_{\mathrm{dd}} = \frac{V}{2}\langle\sigma^{\mathrm{in}}_{jl}\rangle \aleph_{jklm} \langle\sigma^{\mathrm{in}}_{km}\rangle, \tag{5}$$

where the four-rank tensor of "destressing" coefficients $\hat{\aleph}$ depends on the crystal shape and elastic moduli. For example, for a crystal in the form of a thin plate with the normal $\mathbf{n}$, the components of the destressing tensor

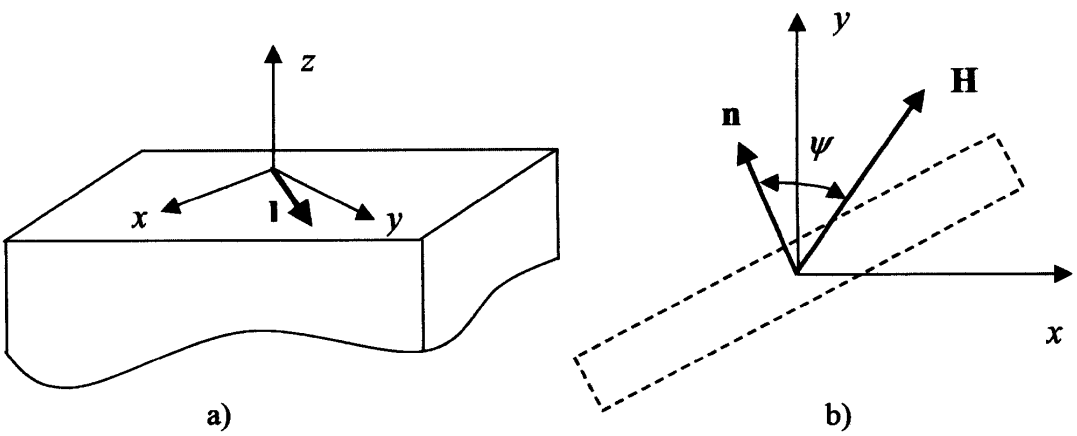

Fig. 1. Crystal shape (a) and orientations of the plate normal and the external magnetic field with respect to the crystallographic axes (b)

are space-independent and may be expressed through the components of the so-called dynamic matrix $D_{jk}^{-1} \equiv c_{ijkl} n_i n_l$ (see, e.g, [23,24]):

$$\aleph_{jklm} = D_{jk} n_l n_m. \tag{6}$$

If, in addition, the plate normal is parallel to one of the so-called pure directions (i.e. if $\mathbf{n}$ is an eigenvector of the $\hat{D}$ matrix), the expression for the destressing energy is reduced to

$$\Phi_{\rm dd} = \frac{V}{2}\left[\frac{1}{c_l}(\mathbf{n}\hat{\sigma}^{(\rm in)}\mathbf{n})^2 + \right.$$

$$\left. + \frac{1}{c_{t1}}(\mathbf{e}_1\hat{\sigma}^{(\rm in)}\mathbf{e}_1)^2 + \frac{1}{c_{t2}}(\mathbf{e}_2\hat{\sigma}^{(\rm in)}\mathbf{e}_2)^2\right], \tag{7}$$

where c_l and $c_{t1,2}$ are the effective moduli corresponding to the longitudinal acoustic mode and two transverse ones propagating in the $\mathbf{n}$ direction.

3. Shape-induced Anisotropy

To elucidate the role of a shape-induced magnetic anisotropy, we consider the typical case of an "easy-plane" AFM with a degenerated equilibrium orientation of the AFM vector $\mathbf{l}$ in the xy plane (bare magnetic in-plane anisotropy is supposed to be vanishingly small, see below). Suppose that the sample is cut in the form of a thin plate oriented perpendicularly to an easy-plane ($\mathbf{n}$ is perpendicular to the z-axis, see Fig. 1,a). Suppose further that the internal stress tensor depends only on the direction of vector $\mathbf{l}$ so that

$$\hat{\sigma}^{(\rm in)} = \lambda_{12}\mathbf{l}^2\hat{1} + \lambda_{44}\mathbf{l}\otimes\mathbf{l}, \tag{8}$$

where λ_{12} and λ_{44} are the isotropic and anisotropic magnetoelastic constants, respectively. The former is responsible for the volume effect that accompanies the magnetic ordering and may include a contribution of the exchange nature, while the latter has mainly the spin-orbit origin.

If the AFM ordering is homogeneous throughout the sample (this may be the case if the characteristic value of the bare in-plane magnetic anisotropy is negligibly small in comparison with effective magnetostrictive fields), then the destressing energy (7) with account of Eq. (8) may be expressed as

$$\Phi_{\rm dd} = V\left[-\frac{1}{2}K_2^{(\rm elas)}(\mathbf{ln})^2 + \frac{1}{4}K_4^{(\rm elas)}(\mathbf{ln})^4\right], \tag{9}$$

where we have introduced the effective constants

$$K_2^{(\rm elas)} = 2\left[\frac{\lambda_{44}^2}{c_t} + \lambda_{12}\lambda_{44}\left(\frac{1}{c_t} - \frac{1}{c_l}\right)\right],$$

$$K_4^{(\rm elas)} = 2\lambda_{44}^2\left(\frac{1}{c_t} + \frac{1}{c_l}\right). \tag{10}$$

Usually, the value of the shear modulus is much less than that of the compression one, $c_t < c_l$. So, effective constant $K_2^{(\rm elas)}$ may be either positive or negative, depending on the sign of $\lambda_{12}\lambda_{44}$, while $K_4^{(\rm elas)} > 0$ is always positive.

It is seen from (9) that the sample shape (described by the direction of the sample normal $\mathbf{n}$) may be a source of the additional anisotropy for the magnetic subsystem whose value is comparable or even much greater than that of the "bare" (proper) magnetic anisotropy. For example, for some crystals (see Table), the "easy-plane"

magnetic anisotropy is below $3 \cdot 10^4$ erg/cm^3, while the shape-induced anisotropy is $K^{(\mathrm{elas})} \approx 10^5$ erg/cm^3.

It should be stressed that the appearance of shape-induced anistropy in AFM is somehow analogous to the effect of induced anisotropy observed in ferromagnetic martensites like Ni-Ga-Mn. In both cases (AFM and FM martensite), the magnetoelastic coupling contributes significantly into the induced anisotropy; in both cases, the internal stresses occur either due to the magnetic (as in AFM) or martensitic (as in FM martensite) phase transition. Meanwhile, the theoretical interpretation of these phenomena is quite different. It is accepted [25, 26] that, in FM martensite, the value of the induced anisotropy is governed mainly by local stresses, while, according to our model, the shape-induced anisotropy in AFM is due to the long-range character of magnetoelastic forces.

Nevertheless, there is no reason to eliminate a possibility for nonlocal magnetoelastic effects to be present in martensites, where spontaneous internal stresses are usually quite large. On the other hand, destressing phenomena in these crystals which possess a nonzero macroscopic magnetization may be veiled by strong demagnetization effects that are also shape-dependent.

To elucidate the role of a sample shape, we neglect the bare in-plane magnetic anisotropy, while the out-plane anisotropy is supposed to be quite large to keep the AFM vector in the "easy plane". Let the magnetic field $\mathbf{H}$ be applied in parallel to the xy plane (see Fig. 1,*b*)). An equilibrium orientation of the AFM vector is then may be calculated by the minimization of expression (3) with a due account of the Zeeman energy in the form (see [27])

$$f^{\mathrm{mag}}[\mathbf{l}(\mathbf{r})] = f_0^{\mathrm{mag}}[\mathbf{l}(\mathbf{r})] - \frac{1}{2}\chi[\mathbf{l} \times \mathbf{H}]^2, \qquad (11)$$

where f_0^{mag} is the magnetic energy density without field, $\chi \equiv M_0/H_E$ is the magnetic susceptibility, and M_0 is the saturation magnetization. In (11), we have taken into account that the external field H is much less than the spin-flip field H_E: $H \ll H_E$.

Depending on the relation between the values of isotropic, λ_{12}, and anisotropic λ_{44} strictions, the minimization of (3),(9), and (11) with respect to the AFM vector components gives rise to the following results. If $\lambda_{12} \gg \lambda_{44}$, then, as seen from Eq. (10), $K_2^{(\mathrm{elas})} \gg K_4^{(\mathrm{elas})}$. In other words, the effective shape-induced anisotropy is equivalent to a biaxial one. If, in addition, $\lambda_{12}\lambda_{44} > 0$, then the "easiest" axis (corresponding to the energetically favourable orientation) is parallel to $\mathbf{n}$. In this case, the angle $\theta = \widehat{(\mathbf{l}, \mathbf{n})}$ between the AFM vector and the plate normal can be calculated from the relation

$$\cos 2\theta = \frac{(H_{\mathrm{eff}}^{(2)})^2 - H^2 + 2[\mathbf{H} \times \mathbf{n}]^2}{\sqrt{[(H_{\mathrm{eff}}^{(2)})^2 - H^2]^2 + 4(H_{\mathrm{eff}}^{(2)})^2[\mathbf{H} \times \mathbf{n}]^2}}, \qquad (12)$$

where we have introduced an effective field $H_{\mathrm{eff}}^{(2)} = \sqrt{K_2^{(\mathrm{elas})}/\chi}$. The analysis of formula (12) shows that, for a small field value $H \ll H_{\mathrm{eff}}^{(2)}$, the AFM vector is almost perpendicular to the plate surface ($\mathbf{l}^{(\mathrm{eq})} \| \mathbf{n}$), while for a relatively large field $H \gg H_{\mathrm{eff}}^{(2)}$, the AFM vector tends to orient perpendicularly to $\mathbf{H}$.

If the magnetostrictive constants λ_{12} and λ_{44} have the same order of magnitude, the effective shape-induced anisotropy is still biaxial, but the easy-axis orientation with respect to the sample surface depends upon a "play of constants". In the particular case of $\lambda_{12} < \lambda_{44} c_t/2c_l$, the equilibrium orientation of the AFM vector in an external field is calculated from the relation

$$(\cos 2\theta_0 - \cos 2\theta)\sin 2\theta - \left(\frac{H}{H_{\mathrm{eff}}^{(4)}}\right)^2 \sin 2(\theta - \psi) = 0,$$

$$\psi = \widehat{(\mathbf{H}, \mathbf{n})}, \qquad (13)$$

where $H_{\mathrm{eff}}^{(4)} = \sqrt{K_4^{(\mathrm{elas})}/2\chi}$ is the effective field, and the angle θ_0 describes the orientation of the in-plane "easy

Comparison of shape-induced $K^{(\mathrm{elas})}$ and bare $K^{(\mathrm{mag})}$ magnetic anisotropies for typical "easy-plane" AFMs. Characteristic values of the specific energy (erg/cm^3) are estimated according to experimental values of the spontaneous magnetostriction $\hat{u}^{(0)}$, saturation magnetization M_0, shear modulus c_t, and spin-flip field H_E

AFM	$K^{(\mathrm{elas})}$	$K^{(\mathrm{mag})}$	$u^{(0)}$	M_0, kGs	c_t, GPa	H_E	Source
$KCoF_3$	$3{\cdot}10^6$	—	$3{\cdot}10^{-3}$	0.5	30	200	[16, 31, 32]
$La_{2-x}Sr_xCuO_4$	$1{\cdot}10^8$	$<5{\cdot}10^6$	$\sim 1{\cdot}10^{-2}$	0.5	100	$\sim$1000	[13, 33, 34]
$CoCl_2$	$5.6{\cdot}10^5$	$< 3{\cdot}10^4$	$4{\cdot}10^{-4}$	0.16	34.7	1.6	[7, 11, 35]
NiO	$0.8{\cdot}10^4$	288	$9{\cdot}10^{-5}$	1.02	109	1000	[29, 36—41]

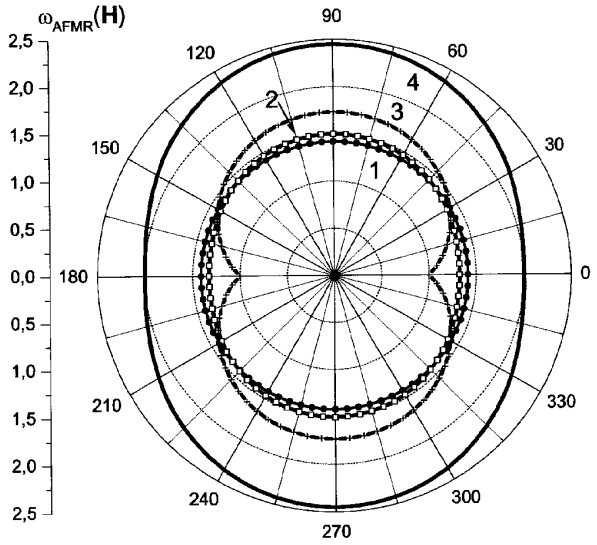

Fig. 2. AFMR frequency gap $\omega_{\rm AFMR}$(dimensionless) vs the angle between the external magnetic field and the plate normal for various field values: *1* — $H = 0,1H_{\rm eff}^{(2)}$, *2* — $H = 0,5H_{\rm eff}^{(2)}$, *3* — $H = H_{\rm eff}^{(2)}$, *4* — $H = 2H_{\rm eff}^{(2)}$; $H_{\rm ms} = H_{\rm eff}^{(2)}$

axis" in the zero field:

$$\cos 2\theta_0 = \frac{2K_2^{(\rm elas)} - K_4^{(\rm elas)}}{K_4^{(\rm elas)}}. \qquad (14)$$

In the limiting case $H \ll H_{\rm eff}^{(4)}$, Eq. (13) has the approximate solution

$$\cos 2\theta = \cos 2\theta_0 + \left(\frac{H}{H_{\rm eff}^{(4)}}\right)^2 (\cos 2\psi - \sin 2\psi \cot 2\theta_0) \qquad (15)$$

that describes the AFM vector rotation under the action of an external field.

For a particular relation between the magnetoelastic constants $\lambda_{12} \approx -\lambda_{44}/2$ (and, respectively, $K_2^{(\rm elas)} \approx K_4^{(\rm elas)}/2$), the effective shape-induced anisotropy may become a uniaxial one. In this case, all the magnetic and magnetoelastic properties of the sample are invariant with respect to the 90° rotation around the z-axis. So, without field, the equilibrium orientation of the AFM vector is twice degenerate in the xy plane, namely, in-parallel or perpendicularly to the plate normal. Correspondingly, the sample should show the spin-flop phenomenon (a turn over of the AFM vector by 90°) at $H = H_{\rm eff}^{(4)}/2$, if the external field is applied in-parallel to one of the "easy" directions.

So, in the crystals with negligibly small magnetic anisotropy (compared with magnetoelastic interactions), the equilibrium orientation of the AFM vector is governed mainly by the sample shape (and may be distinguished from the conventional one). Nevertheless, the shape-induced anisotropy gives rise to the same peculiarities of equilibrium AFM states as a proper magnetic anisotropy. In order to separate the long-range magnetoelastic contribution, one must test the samples with different shapes.

4. Homogeneous AFMR

The magnetic anisotropies of AFM are usually investigated by studying the rotation torques or AFMR spectra in an external magnetic field.

In this section, we consider the effects of the shape-induced anisotropy that may be detected in AFMR spectra.

The AFMR frequency may be easily calculated in the framework of the Lagrange method modified by Bar'yakhtar et al. [27] for the case of AFM. The Lagrangian of AFM with the account of the coupling between the magnetic and elastic subsystems is

$$L = \frac{1}{2}\int\limits_V \left[\frac{\chi}{g^2}\dot{\mathbf{l}}^2 + \rho\dot{\mathbf{u}}^2\right] dV - \Phi. \qquad (16)$$

where g is the gyromagnetic ratio, ρ is the crystal density, and $\mathbf{u}$ is a displacement vector. The last term corresponding to the "potential energy" is given by (3), and we have omitted the inhomogeneous exchange interaction which is immaterial in the long-wave limit.

The analysis of the dynamic equations obtained from the minimization of Lagrangian (16) indicates the existence of two coupled magnetoelastic low-energy modes. One of these modes originates from the activationless transverse acoustic mode. More interesting is another one which corresponds to a low-energy oscillation of the AFM vector in the basal ("easy") plane.

In the case of strong isotropic magnetostriction ($\lambda_{12} \gg \lambda_{44}$, see above), the oscillation frequency is given by the expression

$$\omega_{\rm AFMR}(\mathbf{H}) =$$

$$= g\left\{H_{\rm ms}^2 + \sqrt{[(H_{\rm eff}^{(2)})^2 - H^2]^2 + 4(H_{\rm eff}^{(2)})^2[\mathbf{H}\times\mathbf{n}]^2}\right\}^{1/2}, \qquad (17)$$

where the characteristic field $H_{\rm ms} = \sqrt{2\lambda_{44}^2/\chi c_t}$ defines the isotropic magnetoelastic gap (experimentally observed by Borovik-Romanov et al. [21] and

theoretically explained by Turov et al. [20]). The unconventional feature of the above expression is the *anisotropic* magnetoelastic contribution to the spectrum gap (the last term in (17)) which has the same order of value as the isotropic one. So, the gap in the AFMR spectrum originates from two comparable magnetoelastic contributions. The isotropic term is local and results from the lattice deformation due to the spin occurence at a site, while the anisotropic term arises from nonlocal elastic interactions and is sensitive to the sample shape.

The main peculiarity of frequency (17), namely, the dependence on the mutual orientation of an external field and the plate normal (angle ψ in Fig. 2) is illustrated with Fig. 2. This figure shows the $\omega_{AFMR}(\psi)$ dependence calculated for different external field-to-$H_{\rm eff}^{(2)}$ ratios (for the sake of simplicity, we assume that $H_{\rm ms} = H_{\rm eff}^{(2)}$). At $H \ll H_{\rm eff}^{(2)}$, the energy gap is almost isotropic, while it shows a maximum anisotropy at $H = H_{\rm eff}^{(2)}$. The latter case corresponds to the reorientation of the AFM vector through the "hard" direction (perpendicular to the plate normal). With a further increase in the field H, the gap anisotropy becomes less pronounced. The maximum/ minimum frequency $\omega_{\rm AFMR}^{\rm max,min} = g\sqrt{H_{\rm ms}^2 + |(H_{\rm eff}^{(2)})^2 \pm H^2|}$ is attained at $\mathbf{H} \perp \mathbf{n}$ and $\mathbf{H} \| \mathbf{n}$, respectively.

The above-mentioned peculiarities of AFMR spectra may be observed for a wide range of the magnetostrictive constants $\lambda_{12}, \lambda_{44}$ and their ratio. Some other behavior may be expected in a particular case $\lambda_{12} \approx -\lambda_{44}/2$, when the ground state of AFM (without field) is doubly degenerate. At that, the AFMR spectrum consists of two branches $\omega_{AFMR}^{\pm}$ that correspond to the ground states with $\mathbf{l} \| \mathbf{n}$ and $\mathbf{l} \perp \mathbf{n}$. For $H \ll H_{\rm eff}^{(4)}$,

$$\omega_{\rm AFMR}^{\pm}(\mathbf{H}) = g\sqrt{H_{\rm ms}^2 + \frac{1}{2}(H_{\rm eff}^{(4)})^2 \mp 2H^2 \pm 4[\mathbf{H} \times \mathbf{n}]^2}. \tag{18}$$

So, an external field removes the degeneracy of the ground states with $\mathbf{l} \| \mathbf{n}$ and $\mathbf{l} \perp \mathbf{n}$. From the energetic point of view, the state for which the AFM vector $\mathbf{l}$ makes a larger angle with an external field is more preferable (the corresponding frequency $\omega_{\rm AFMR}^{+}(\mathbf{H})$). The experimental observation of one or another branch depends upon the prehistory of a sample and the kinetics of relaxation to an equilibrium state after the application of the external magnetic field.

5. Conclusions

In conclusion, the crystal shape greatly affects the equilibrium orientation of the AFM vector. In the crystals with strong magnetoelastic coupling, the shape is a source of the additional (as compared to bare magnetic) anisotropy and establishes some new "easy" directions, whose orientation is governed by the external properties of a sample rather than the internal ones. The shape effect results from the long-range interaction of internal stresses induced by the AFM ordering and rigidly coupled with a local orientation of the AFM vector.

The shape-induced anisotropy may be detected experimentally by studying the rotation torques or AFMR spectra in an external magnetic field. For a thin plate, the value of a rotation torque and the frequency gap in the spin-wave spectrum should depend upon the mutual orientation of the plate normal and the direction of the external magnetic field as $[\mathbf{H} \times \mathbf{n}]^2$. This dependence should reveal itself even for the samples isotropic from the crystallographic point of view.

Up to our knowledge, there were no special experiments aimed at the investigation of shape-induced effects in AFM. However, Roth and Slack [28] noticed a strong and nontrivial shape-dependence of the magnetic anisotropy which revealed itself in the torque measurements of AFM NiO. In particular, the rotational torques when the field is rotated in (111) plane show the 90°-periodicity vs the magnetic field direction for the samples with square cross-section, while the crystallographic magnetic anisotropy requires the 60°-periodicity [29,30]. It can be easily shown that such a shape-related change of the effective anisotropy is in accordance with the above model.

It should be noted that the shape-induced or proper magnetic anisotropy may give rise to similar effects. In order to ascertain the role and importance of the shape-induced anisotropy and distinguish it from the crystallographic one, it is necessary to test the samples of different shapes, for example, ellipsoids with different aspect ratio. The theoretical analysis of different shapes is the subject for further investigations.

We are grateful to V.A. L'vov for the informal review of the paper and the constructive criticism. H.V.G. acknowledges the technical and financial support from A.A. Malyshenko.

This study was done as a part of the project financed by the Ministry of Education of Ukraine (State Registration No. 0105U001280).

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ПРО МАГНІТОПРУЖНУ ВЗАЄМОДІЮ ЯК ПРИЧИНУ ВПЛИВУ ФОРМИ КРИСТАЛА НА СПЕКТР АФМР

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Р е з ю м е

На основі феноменологічної моделі з урахуванням енергії роздеформування показано, що в антиферомагнітних кристалах із сильним магнітопружним зв'язком форма зразка приводить до виникнення "штучної" анізотропії магнітопружної природи, величина якої може перебільшувати власну магнітну анізотропію. Така наведена формою анізотропія повинна приводити до залежності щілини в спектрі АФМР від напрямку зовнішнього магнітного поля.

О МАГНИТОУПРУГОМ ВЗАИМОДЕЙСТВИИ КАК ПРИЧИНЕ ВЛИЯНИЯ ФОРМЫ КРИСТАЛЛА НА СПЕКТР АФМР

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Р е з ю м е

В рамках феноменологической модели с учетом энергии раздеформирования показано, что в антиферромагнитных кристаллах с сильной магнитоупругой связью форма образца приводит к возникновению "искусственной" анизотропии магнитоупругой природы, величина которой может превышать собственную магнитную анизотропию. Такая наведенная формой анизотропия должна приводить к зависимости щели в спектре АФМР от направления внешнего магнитного поля.