Domain-wall drift in antiferromagnets

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A theory is developed describing the drift of a planar 180-degree domain wall in a collinear antiferromagnet with equivalent sublattices in a weak external magnetic field oscillating with an arbitrary frequency. The character of the dependence of the drift velocity on the frequency and polarization of the oscillating field is determined for different ratios of the parameters of the uniaxial and planar anisotropy.

A reasonably complete theory of the drift motion of domain walls (DWs) or their elements in an oscillating field has been constructed for ferromagnets (FMs). This was made possible by the simplicity of the single-sublattice model. There exist several approaches to solving the problem of the dynamic response of such a magnet to nonlinear external excitations. According to Schlömann,[1] in a circularly polarized field a domain wall should be subjected to an effective pressure, arising due to the difference of the magnetic energy density on both sides of the wall. Another approach employs the assumption that the frequency ω of the oscillating field is much higher than the frequency of the characteristic oscillations of the magnet[2,3] or the frequency of oscillations of the domain wall.[4] Finally, a systematic theory was recently developed for domain-wall drift without the restriction $\omega \gg \omega_r$ but assuming that the amplitude of the oscillating field is small.[5,6] These theories are supported by many experimental results, most of which were obtained in our country.

At the same time, domain-wall motion in an external oscillating or nonuniform magnetic field is characteristic not only for magnets with ferromagnetic ordering. In particular, Bar'yakhtar *et al.*[7] pointed out that such motion is possible in antiferromagnets (AFMs), and this has been established experimentally.[8]

In the present paper a systematic asymptotic theory is constructed for the drift of a flat domain wall in a multisublattice magnet in a weak oscillating field. The oscillating field is assumed to be elliptically polarized, and the limiting cases of high and low frequencies of the field are studied.

1. EFFECTIVE EQUATIONS OF MOTION OF THE MAGNETIZATION OF AN AFM IN AN OSCILLATING FIELD

In order to describe the dynamics of a two-sublattice anisotropic AFM we start from the equations for the normalized ferromagnetism vector **m** and antiferromagnetism vector **l**, which are related to the magnetization vectors $\mathbf{M}_1$ and $\mathbf{M}_2$ of the sublattices by the relations

$$\mathbf{m}=\frac{\mathbf{M}_1+\mathbf{M}_2}{2M_0},\quad \mathbf{l}=\frac{\mathbf{M}_1-\mathbf{M}_2}{2M_0}. \tag{1}$$

Setting $|\mathbf{M}_1| = |\mathbf{M}_2| = M_0$ we obtain the conditions which relate the vectors **m** and **l**:

$$\mathbf{m}^2+\mathbf{l}^2=1,\quad (\mathbf{ml})=0. \tag{2}$$

The dynamics of these vectors in the presence of dissipation is described by the system of Landau–Lifshitz equations:

$$\begin{aligned}\dot{\mathbf{m}}&=-\frac{\gamma}{2}([\mathbf{mH}_m]+[\mathbf{lH}_l])+\frac{\varepsilon}{M_0}([\mathbf{m}\dot{\mathbf{m}}]+[\mathbf{l}\dot{\mathbf{l}}]),\\ \dot{\mathbf{l}}&=-\frac{\gamma}{2}([\mathbf{mH}_l]+[\mathbf{lH}_m])+\frac{\varepsilon}{M_0}([\mathbf{m}\dot{\mathbf{l}}]+[\mathbf{l}\dot{\mathbf{m}}]),\end{aligned} \tag{3}$$

where γ is the gyromagnetic ratio, ε is the damping parameter,

$$M_0\mathbf{H}_m=-\frac{\delta W}{\delta\mathbf{m}},\quad M_0\mathbf{H}_l=-\frac{\delta W}{\delta\mathbf{l}}$$

are the effective magnetic fields, and W is the energy of the AFM.

We represent the energy of a collinear AFM with orthorhombic symmetry in an alternating field as a sum of the energy W_0 of the AFM itself and the energy of the alternating field W_h:

$$W=W_0+W_h, \tag{4}$$

$$W_0=M_0^2\int_{-\infty}^{\infty} d\mathbf{r}\{2\delta\mathbf{m}^2+\alpha_1(\nabla\mathbf{l})^2+\alpha_2(\nabla\mathbf{m})^2-W_a-2(\mathbf{mh}_C)\}, \tag{5}$$

$$W_h=M_0^2\int_{-\infty}^{\infty} d\mathbf{r}\{-2(\mathbf{mh}^{(\omega)})\},$$

$$W_a=\beta_1 l_z^2+\beta_2 m_z^2+\rho_1 l_y^2+\rho_2 m_y^2. \tag{6}$$

Here δ, α_1, and α_2 are the homogeneous and inhomogeneous exchange constants, respectively; β_1, β_2 and ρ_1, ρ_2 are the uniaxial and planar anisotropy constants; $\mathbf{h}_C = \mathbf{H}_C/M_0$ is the constant external field, which we assume is parallel to the uniaxial anisotropy axis $\mathbf{H}_C = \mathbf{H}_C\mathbf{e}_z$; $\mathbf{h}^{(\omega)} = \mathbf{H}^{(\omega)}/M_0$ is the oscillating magnetic field, where $\mathbf{H}^{(\omega)} = \mathrm{Re}(\tilde{\mathbf{H}}\exp(i\omega t))$.

We determine the ratio of the anisotropy constants in a manner so that the domain wall is oriented in the zx plane and the ferro- and antiferromagnetism vectors rotate in the plane of the wall

$$\beta_1-\beta_2>0,\quad \rho_1-\rho_2<0.$$

We substitute the explicit form of the effective fields into Eqs. (3) and then take into account the fact that in the long-wavelength approximation all terms containing gradients of the magnetization, together with the terms stemming from the magnetic anisotropy energy, can be dropped.[9] As a

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result, since the inequality $\mathbf{m}^2 \ll \mathbf{l}^2$ is also satisfied, as is obvious for a collinear AFM, we obtain the following equations for the vectors **m** and **l**:

$$\mathbf{m}=\frac{1}{2\delta\gamma M_0}[\mathbf{l}\dot{\mathbf{l}}]+\frac{1}{2\delta M_0}\{\mathbf{H}-\mathbf{l}(\mathbf{H}\mathbf{l})\}, \tag{7}$$

$$[\mathbf{l}\ddot{\mathbf{l}}]-c^2[\mathbf{l}\mathbf{l}'']=\omega_0^2 l_z[\mathbf{l}\mathbf{e}_z]+\omega_1^2 l_y[\mathbf{l}\tilde{\mathbf{e}}_y]-\nu[\mathbf{l}\dot{\mathbf{l}}]$$
$$-\gamma^2(\mathbf{l}\mathbf{H})[\mathbf{l}\mathbf{H}]-\gamma\{2\dot{\mathbf{l}}(\mathbf{l}\mathbf{H})-\dot{\mathbf{H}}^{(\omega)}+\mathbf{l}(\mathbf{l}\dot{\mathbf{H}}^{(\omega)})\}, \tag{8}$$

$$\mathbf{H}=\mathbf{H}_c+\mathbf{H}^{(\omega)},\quad \omega_0^2=2\delta\beta_1(\gamma M_0)^2,\quad \omega_1^2=2\delta\rho_1(\gamma M_0)^2,\quad \nu=2\varepsilon\delta\gamma,$$
$$c^2=2\delta\alpha_1(\gamma M_0)^2, \tag{9}$$

where the prime denotes differentiation with respect to y; c is the characteristic velocity, which for $\mathbf{H}=0$ is equal to the minimum phase velocity of spin waves in the linear theory; and, ω_0 is the frequency of uniform resonance of magnetization oscillations with $\mathbf{H}=0$ (Ref. 9). It is easy to see that on the one hand the relations obtained above with $\mathbf{H}^{(\omega)}=0$ are identical to the analogous well-known formulas of Ref. 9, while on the other hand they show that they can also be used to describe nonlinear waves of magnetization of an AFM in the absence of a constant field, the role of which, in the last case, the field $\mathbf{H}^{(\omega)}$ plays. For this reason, in what follows the constant field is not necessarily taken into account in Eq. (8).

2. PERTURBATION THEORY

By analogy to ferromagnets[6] we assume that the dynamic distribution of the vector **1** in the domain wall of an antiferromagnet can be represented in the form

$$\mathbf{l}(y,t)=\mathbf{L}(\tilde{y})+\mathbf{n}(\tilde{y},t), \tag{10}$$

where $y=y-Y(t)$, $Y(t)$ has the meaning of a coordinate, and $\mathbf{L}(y)$ and $\mathbf{n}(y,t)$ describe the distribution $\mathbf{l}(y,t)$ in the coordinate system of the domain wall.

Substituting Eq. (10) into Eq. (8) we obtain

$$[\mathbf{L}\ddot{\mathbf{n}}]-\ddot{Y}[\mathbf{L}\mathbf{n}']+(\dot{Y})^2[\mathbf{L},\mathbf{L}''+\mathbf{n}'']-\dot{Y}[\mathbf{L},\mathbf{L}'+\mathbf{n}']-c^2[\mathbf{L}\mathbf{L}'']$$
$$-\omega_0^2 L_z[\mathbf{L}\mathbf{e}_z]-\omega_1^2 L_y[\mathbf{L}\mathbf{e}_y]+\nu[\mathbf{L},\dot{\mathbf{n}}-\dot{Y}(\mathbf{n}'+\mathbf{L}')]$$
$$+\gamma^2(\mathbf{L}\mathbf{H}^{(\omega)})[\mathbf{L}\mathbf{H}^{(\omega)}]$$
$$+\gamma\{2(\mathbf{L}\mathbf{H}^{(\omega)})(\dot{\mathbf{n}}-\dot{Y}(\mathbf{n}'+\mathbf{L}'))-\dot{\mathbf{H}}^{(\omega)}+\mathbf{L}(\mathbf{L}\dot{\mathbf{H}}^{(\omega)})\}=0. \tag{11}$$

Next we set

$$\mathbf{L}(y)=\mathbf{L}_0(y)+\mathbf{L}_1(y)+\mathbf{L}_2(y)+\ldots,$$
$$\mathbf{n}(y,t)=\mathbf{n}_1(y,t)+\mathbf{n}_2(y,t)+\ldots, \tag{12}$$
$$\dot{Y}(t)=V_1+V_2+\ldots+\dot{u}_1(t)+\dot{u}_2(t)+\ldots,$$
$$\ddot{Y}(t)=\dot{V}_1+\dot{V}_2+\ldots+\ddot{u}_1(t)+\ddot{u}_2(t)+\ldots$$

with zero mean values of the oscillating variables

$$\mathbf{n}_i(y,t)\ \text{and}\ u_i(t)$$
$$\langle\mathbf{n}_i\rangle=\langle\dot{u}_i\rangle=\langle\ddot{u}_i\rangle=0,\quad i=1,2,\ldots, \tag{13}$$

where the numerical indices indicate the order of smallness of the quantities relative to the amplitude of the oscillating field.

Since we are interested in the stationary motion of a domain wall, we write out, using Eqs. (12) and (13), the zeroth (ground state) and first two orders of the perturbation theory for Eq. (11):

$$c^2[\mathbf{L}_0\mathbf{L}_0'']+\omega_0^2(L_0)_z[\mathbf{L}_0\mathbf{e}_z]+\omega_1^2(L_0)_y[\mathbf{L}_0\mathbf{e}_y]=0, \tag{14}$$

$$[\mathbf{L}_0\ddot{\mathbf{n}}_1]-\ddot{u}_1[\mathbf{L}_0\mathbf{L}_0']-c^2([\mathbf{L}_0\mathbf{n}_1'']+[\mathbf{n}_1\mathbf{L}_0''])$$
$$-\omega_0^2((L_0)_z[\mathbf{n}_1\mathbf{e}_z]+(n_1)_z[\mathbf{L}_0\mathbf{e}_z])$$
$$-\omega_1^2((L_0)_y[\mathbf{n}_1\mathbf{e}_y]+(n_1)_y[\mathbf{L}_0\mathbf{e}_y])$$
$$+\nu([\mathbf{L}_0\dot{\mathbf{n}}_1]-\dot{u}_1[\mathbf{L}_0\mathbf{L}_0'])+\gamma\{\mathbf{L}_0(\mathbf{L}_0\dot{\mathbf{H}}^{(\omega)})-\dot{\mathbf{H}}^{(\omega)}\}=0, \tag{15}$$

$$\langle[\mathbf{n}_1\ddot{\mathbf{n}}_1]\rangle-\langle\ddot{u}_1[\mathbf{L}_0\mathbf{n}_1']\rangle+\langle(\dot{u}_1)^2\rangle[\mathbf{L}_0\mathbf{L}_0'']$$
$$-\langle\ddot{u}_1(2[\mathbf{L}_0\mathbf{n}_1']+[\mathbf{n}_1\mathbf{L}_0'])\rangle$$
$$-c^2([\mathbf{L}_0\mathbf{L}_2'']+[\mathbf{L}_2\mathbf{L}_0'']+\langle[\mathbf{n}_1\mathbf{n}_1'']\rangle)$$
$$-\omega_0^2((L_0)_z[\mathbf{L}_2\mathbf{e}_z]+(L_2)_z[\mathbf{L}_0\mathbf{e}_z]$$
$$+\langle(n_1)_z[\mathbf{n}_1\mathbf{e}_z]\rangle)-\omega_1^2((L_0)_y[\mathbf{L}_2\mathbf{e}_y]$$
$$+(L_2)_y[\mathbf{L}_0\mathbf{e}_y]+\langle(n_1)_y[\mathbf{n}_1\mathbf{e}_y]\rangle)$$
$$-\nu\{V_2[\mathbf{L}_0\mathbf{L}_0']+2\langle\dot{u}_1[\mathbf{L}_0\mathbf{n}_1']\rangle-\langle[\mathbf{n}_1\dot{\mathbf{n}}_1]\rangle+\langle\dot{u}_1[\mathbf{n}_1\mathbf{L}_0']\rangle\}$$
$$+\gamma^2\langle(\mathbf{L}_0\mathbf{H}^{(\omega)})[\mathbf{L}_0\mathbf{H}^{(\omega)}]\rangle+\gamma\{\langle(\dot{\mathbf{n}}_1-\dot{u}_1\mathbf{L}_0')(\mathbf{L}_0\mathbf{H}^{(\omega)})\rangle$$
$$+\langle\mathbf{L}_0(\mathbf{n}_1\dot{\mathbf{H}}^{(\omega)})\rangle+\langle\mathbf{n}_1(\mathbf{L}_0\dot{\mathbf{H}}^{(\omega)})\rangle\}=0. \tag{16}$$

These same conditions (12) and (13) applied to the relations (2) yield

$$\mathbf{L}_0^2=1,\quad \mathbf{L}_0\perp\mathbf{n}_1,\quad \langle(\mathbf{n}_1)^2\rangle=-2\mathbf{L}_0\mathbf{L}_2. \tag{17}$$

The equation of the ground state (14) has a well-known solution in the form of a quasi-Bloch planar domain wall

$$\mathbf{L}_0(y)=\mathbf{e}_x\sin\theta(y)+\mathbf{e}_z\cos\theta(y), \tag{18}$$
$$\sin\theta(y)=\mathrm{ch}^{-1}(y/\Delta),\quad \cos\theta(y)=-\mathrm{th}(y/\Delta),\quad \Delta=c/\omega_0.$$

The equations of the higher order aproximations (15) and (16) are most simply studied in a local Cartesian coordinate system $\{\mathbf{e}_\perp,\mathbf{e}_y,\mathbf{e}_\parallel\}$, which we orient so that

$$\mathbf{e}_\parallel\equiv\mathbf{L}_0(y),\quad \mathbf{n}_1=\mathbf{e}_\perp(n_1)_\perp+\mathbf{e}_y(n_1)_y,$$
$$\mathbf{L}_2=\mathbf{e}_\perp(L_2)_\perp+\mathbf{e}_y(L_2)_y+\mathbf{e}_\parallel(L_2)_\parallel. \tag{19}$$

In what follows we drop the order indices of the quantities $\mathbf{n}_1$ and $\mathbf{L}_2$.

As a result the equations for both the oscillating variables (15) and the average variables (16) split into independent equations for the components $\tilde{n}_\perp(\xi)$, $\tilde{n}_y(\xi)$ and $L_\perp(\xi)$, $L_y(\xi)$:

$$\left[\hat{\mathcal{H}}+\left(\frac{\omega}{\omega_0}\right)^2\left(1-\frac{i\nu}{\omega}\right)\right]\tilde{n}_\perp(\xi)$$
$$=\left(\frac{\omega}{\omega_0}\right)^2\left[\frac{\tilde{u}_1}{\Delta}\left(1-\frac{i\nu}{\omega}\right)\sin\theta-\frac{i\gamma}{\omega}\tilde{H}_\nu^{(\omega)}\right], \tag{20}$$

$$\left[\hat{\mathscr{H}}-\left(\frac{\omega}{\omega_0}\right)^2\left(\left(1-\frac{i\nu}{\omega}\right)+\left(\frac{\omega_1}{\omega}\right)^2\right)\right]\tilde{n}_y(\xi)$$

$$=\left(\frac{\omega}{\omega_0}\right)^2\frac{i\gamma}{\omega}\left(\tilde{H}_x^{(\omega)}\cos\theta-\tilde{H}_z^{(\omega)}\sin\theta\right), \tag{21}$$

$$\hat{\mathscr{H}}L_\perp(\xi)=\langle N\rangle,$$

$$N=(nn'-n_\perp n_\perp')2\sin\theta-(n_\perp)^2\sin2\theta$$

$$-\left(\frac{1}{\omega_0^2}\right)\Big[\dot{u}_1\dot{n}_\perp'+2(\ddot{u}_1 n_\perp')$$

$$-\left(\frac{\dot{u}_1}{\Delta}\right)^2\sin\theta\cos\theta+\nu\left(\frac{V_2}{\Delta}\sin\theta+2(\dot{u}_1 n_\perp')\right)\Big]$$

$$+\left(\frac{\gamma}{\omega_0}\right)^2[H_x^{(\omega)}H_z^{(\omega)}\cos2\theta+((H_x^{(\omega)})^2-(H_z^{(\omega)})^2)\sin\theta\cos\theta]$$

$$+\left(\frac{\gamma}{\omega_0^2}\right)\frac{d}{dt}[n_y(H_x^{(\omega)}\sin\theta+H_z^{(\omega)}\cos\theta)]. \tag{22}$$

Here

$$\hat{\mathscr{H}}=\frac{d^2}{d\xi^2}-\cos2\theta(\xi),\quad \xi=\frac{y}{\Delta},$$

$$\mathbf{n}(\xi,t)=\mathrm{Re}\,[\tilde{\mathbf{n}}(\xi)\exp(i\omega t)],\quad u_1(t)=\mathrm{Re}\,[\tilde{u}\exp(i\omega t)].$$

We do not write out the equation for $L_y(\xi)$, since it does not contain the desired quantity V_2 and for this reason it is of no further interest to us.

The obtained equations include the negative-definite self-adjoint operator $\hat{\mathscr{H}}$. The zero eigenvalue of this operator, describing a uniform shift of the domain wall, corresponds to the eigenfunction $\theta'(\xi)=\sin\theta(\xi)$. Hence it is easy to find the solutions of these equations. Thus the condition for Eq. (22) to have a solution is that the right-hand side of the equation be orthogonal to the eigenfunction of the operator $\hat{\mathscr{H}}$. As a result, we obtain an expression for the drift velocity of the domain wall similar to the corresponding expression for a ferromagnet:[6]

$$V_2=\frac{\Delta\omega_0^2}{2\nu}\int_{-\infty}^{\infty}\theta'(\xi)\langle N\rangle d\xi. \tag{23}$$

3. APPLICATION OF THE THEORY

Let the oscillating field be polarized in the plane of the domain wall, i.e., $\tilde{H}_y^{(\omega)}=0$. In this case the solution of Eq. (20) can be found in the form $\tilde{n}_\perp(\xi)=a_1\sin\theta+a_2\cos\theta$. Substituting it into Eq. (20) we find

$$a_1=\tilde{u}_1/\Delta,\quad a_2=0. \tag{24}$$

The solution of Eq. (21) can be found similarly, and it is equal to

$$\tilde{n}_y(\xi)=\frac{i\omega\gamma\tilde{H}_z^{(\omega)}}{\omega(\omega+i\nu)+\omega_1^2}\sin\theta-\frac{i\omega\gamma\tilde{H}_x^{(\omega)}}{\omega(\omega+i\nu)+\omega_1^2+\omega_0^2}\cos\theta. \tag{25}$$

Using the solutions (24) and (25), we obtain from Eq. (23) an explicit expression for V_2:

$$V_2=V_d\,\mathrm{Re}(S), \tag{26}$$

where

$$S=\left[\left(\frac{\omega^2+\omega_1^2}{\omega\omega_0}\right)^2+\left(\frac{\omega_1}{\omega}\right)^2+\left(\frac{\nu}{\omega_0}\right)^2+1-\frac{i\nu}{\omega}\right]^{-1}(\tilde{h}^{(\omega)})^2, \tag{27}$$

$$V_d=\frac{\pi}{8}\frac{\Delta(\gamma M_0)^2}{\nu},\quad (\tilde{h}^{(\omega)})^2=\frac{\tilde{H}_z^{(\omega)}(\tilde{H}_x^{(\omega)})^*}{M_0^2}. \tag{28}$$

We shall analyze the limiting cases.

First we note that there is a definite relation between the drift velocity, and the symmetry of the AFM. In this connection we distinguish AFMs having orthorhombic symmetry ($\omega_0\approx\omega_1$), substantially uniaxial AFMS ($\omega_0>\omega_1\gtrsim\nu$), and uniaxial AFMs ($\omega_1=0$). The parameter ν, determined by relaxation and uniform exchange, also plays a characteristic role.

For orthorhombic AFMs the ratio of the parameter ν and the frequency of the oscillating field ω is of no fundamental significance, if they are significantly lower than the uniform-resonance frequency ω_0 ($\nu,\omega\ll\omega_0$):

$$V_2=\frac{V_d}{2}\left(\frac{\omega}{\omega_1}\right)^2\mathrm{Re}(\tilde{h}^{(\omega)})^2. \tag{29}$$

For a substantially uniaxial AFM, with the same restriction on the values of ν and ω depending oin their ratio, we have

$$V_2=V_d\left(\frac{\omega}{\omega_1}\right)^2\mathrm{Re}(\tilde{h}^{(\omega)})^2\quad\text{for}\quad\omega\ll\nu, \tag{30}$$

$$V_2=\frac{V_d}{2}\mathrm{Re}(\tilde{h}^{(\omega)})^2\quad\text{for}\quad\omega\gg\nu. \tag{31}$$

The drift velocity in a purely uniaxial AFM is equal to

$$V_2=V_d\left(\frac{\omega}{\nu}\right)\mathrm{Im}(\tilde{h}^{(\omega)})^2\quad\text{for}\quad\omega\ll\nu, \tag{32}$$

$$V_2=V_d\,\mathrm{Re}(\tilde{h}^{(\omega)})^2\quad\text{for}\quad\omega\gg\nu. \tag{33}$$

It is obvious from the expressions (30) and (32) that for low frequencies of the oscillating field ($\omega\ll\nu\ll\omega_0$) the drift velocity is proportional to the small quantities (ω/ν) and $(\omega/\omega_1)^2$, and since ν and ω_1 are fixed parameters, it is obvious that the drift velocity decreases as the frequency ω decreases. For orthorhombic AFMs this result is also valid at frequencies $\omega\gg\nu$ (29). For substantially uniaxial and purely uniaxial AFMs in the frequency interval $\nu\ll\omega\ll\omega_0$ the drift velocity does not depend on the frequency of the oscillating field and is proportional only to the squared amplitude of the field. The strongest effect should be expected in this case [formulas (31) and (33)].

If the frequency of the oscillating field is much higher than the frequency of uniform resonance ($\omega\gg\omega_0$), then the drift velocity is determined uniquely and is equal to

$$V_2=V_d\left(\frac{\omega_0}{\omega}\right)^2\mathrm{Re}(\tilde{h}^{(\omega)})^2. \tag{34}$$

This result agrees with the well-known results of Refs. 2–4, where it is shown that the effect of a rapidly oscillating field is proportional to the parameter $(\gamma \widetilde{H}^{(\omega)}/\omega)^2$.

Substituting typical values for AFMs $\Delta \sim 10^{-5}$ cm, $\delta \sim 10^3$, $\beta \sim \rho \sim 1$, $M_0 \sim 10^2$ Oe, $\lambda \sim 10^{-4}$ ($\lambda = \varepsilon/M_0$ is the dimensionless relaxation constant), we obtain $\omega_0 \sim 10^{11}\ \mathrm{s}^{-1}$, $\nu \sim 10^8\ \mathrm{s}^{-1}$, $V_d \sim 10^5$ m/s.

The amplitude of the oscillating field (the small parameter of the problem) is a quite variable parameter. The point is that although the proposed theory is applicable if the amplitude is small, the upper limit for the theory is still determined by the reasonable requirement that the drift velocity should not exceed the velocity of the domain wall for the given magnet in an external oscillating field.

From the limiting relations presented above [especially the relations (31) and (33)] it is obvious that even for $\widetilde{H} < 1$ the drift velocity of the domain wall has a significant value:

$$V_2 \sim 10^5 (\tilde{h}^{(\omega)})^2 \text{ m/s}. \tag{35}$$

An appreciable effect can also be observed in all other cases with the appropriate choice of the parameters of the oscillating field.

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