Spin-flip transitions in magnets with anisotropy-energy fluctuations

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We consider the features of spin-flip transitions in systems with fluctuations of the magnetic-anisotropy energy and of the antisymmetric exchange. We show that the anisotropy-energy fluctuations increase the width of the spin-flip transition and lead to anomalies of the temperature dependence of the angles that specify the orientation of the average magnetic moment of the system. In noncollinear antiferromagnets with fluctuations of the antisymmetric exchange, it is noted that a spin-flip transition can be induced by an external field applied parallel to the average magnetic moment. A phase diagram in coordinates (T, H) is constructed with account taken of the reversal of the sign of the first anisotropy constant.

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1. INTRODUCTION

In the theoretical study of the properties of disordered magnets, hardly any attention was paid to problems connected with spin-flip transitions in these compounds. Yet recent[1-3] investigations of the disordered substituted ortho-ferrites $RFe_{1-x}M_xO_3$ (R is a rare-earth or Y ion, and M is an iron-group ion), which are classical objects for the study of spin-flip transitions, have revealed a number of anomalies which have received no adequate theoretical explanation to this day. In this paper we analyze the possible features of spin flip in disordered systems with fluctuations of the energy of the anisotropic interactions.

It is known[4,5] that the magnetic structure of disordered magnets can be described by specifying the statistical characteristics of the random field of the vectors $\mathbf{M}_v(\mathbf{r})$, the magnetic moments of the sublattices numbered v. Of particular interest here are the average moments $\langle \mathbf{M}_v \rangle$. The reason is that they determine macroscopic quantities, such as the magnetic moment of the system, and can be investigated by traditional methods of spin-flip observation, whereas to obtain information on the statistical characteristics of the field of the fluctuations $\delta \mathbf{M}_v(\mathbf{r}) = \mathbf{M}_v(\mathbf{r}) - \langle \mathbf{M}_v \rangle$ special experimental procedures are required.[4,5] It is therefore desirable to exclude somehow the fluctuations $\delta \mathbf{M}_v$ from the thermodynamic potential of the system and to use in the analysis of spin-flip transitions the obtained effective potential that depends only on $\langle \mathbf{M}_v \rangle$. In the first section of the paper this approach is applied to systems with random competing magnetic anisotropy. In the second section we consider systems with fluctuations of the antisymmetric Dzyaoshinskiĭ-Moriya exchange,[6] which exerts a substantial influence on the magnetic properties of noncollinear ferromagnets of the $RFeO_3$ type.[7]

2. SPIN-FLIP TRANSITIONS IN SYSTEMS WITH FLUCTUATIONS OF THE MAGNETIC-ANISOTROPY ENERGY

The thermodynamic-potential density in the systems considered, which are characterized by inhomogeneity of the magnetic-anisotropy energy and by the ensuing fluctuations of the vectors $\mathbf{M}_v(\mathbf{r})$, can be represented in the form

$$\Omega = \lim_{V \to \infty} V^{-1} \int \{\omega_{\mathrm{ex}} + \omega_{\mathrm{an}} + \omega_H\} dV, \qquad (1)$$

where V is the volume of the system, and $\omega_{\mathrm{ex}}, \omega_{\mathrm{an}}, \omega_H$ corresponds to the energies of the exchange, anisotropy, and interaction with the external field. We confine ourselves hereafter to the most common transitions, in which the system magnetic moment does not leave a certain selected plane. We put accordingly

$$\omega_{an} = -\frac{1}{2}\beta(\mathbf{r}) \cos^2 \theta(\mathbf{r}), \qquad (2)$$

where $\theta(\mathbf{r})$ is the angle that specifies the magnetic-moment orientation relative to the anisotropy axis. To highlight the effects inherent in disordered systems, we shall disregard higher-order anisotropy. The characteristics of the random field $\beta(\mathbf{r})$ are the mean value $\overline{\beta}$ and the correlation function, which we choose in the form

$$K_\beta(\mathbf{r}-\mathbf{r}') = \langle \beta(\mathbf{r})\beta(\mathbf{r}') \rangle - \overline{\beta}^2 = D(\beta) \exp(-|\mathbf{r}-\mathbf{r}'|/R_c), \qquad (3)$$

where $D(\beta)$ is the dispersion and R_c is the correlation radius of the anisotropy fluctuations. The symbol $\langle \ \rangle$ is used here and elsewhere to denote averaging over the fluctuations of the principal interactions that determine the magnetic structure. The exchange energy is of the usual form:

$$\omega_{\mathrm{ex}} = \frac{1}{2}\alpha(\nabla\theta)^2, \qquad (4)$$

where α is a nonrandom exchange parameter.[1]

We confine ourselves in this section to spontaneous transitions and put accordingly $\omega_H = 0$. Minimizing (1) with respect to θ, we get

$$\alpha\Delta\theta - \beta(\mathbf{r}) \sin\theta \cos\theta = 0. \qquad (5)$$

We assume that the spin-flip fluctuations $\delta\theta(\mathbf{r}) = \theta(\mathbf{r}) - \langle \theta \rangle$ connected with the fluctuations $\delta\beta(\mathbf{r}) = \beta(\mathbf{r}) - \overline{\beta}$ of the energy of the magnetic anisotropy are small. In addition, we neglect in (5) the term $\delta\beta\delta\theta$. These two assumptions are basic in this paper. In this case the angle $\langle \theta \rangle \equiv \overline{\theta}$ determines the orientation of the average magnetic moment of the system. Linearizing (5) with allowance for the assumptions made and taking the Fourier transform, we obtain an equation for the components $\delta\theta(\mathbf{k})$ in the form

$$\delta\theta(\mathbf{k}) = \frac{\delta\beta(\mathbf{k})\sin\theta\cos\theta}{[\alpha k^2 + a(\theta)]}, \qquad a(\theta) = \beta\cos 2\theta. \qquad (6)$$

Using (6) we can determine the basic statistical characteristics of the random field of the fluctuations $\delta\theta(\mathbf{k})$. In particular, the correlation function takes the form

$$K_\theta(\mathbf{r}_1\mathbf{r}_2) = K_\theta(\mathbf{r}_{12})$$

$$= \frac{\sin^2\theta\cos^2\theta}{a^{1/2}(\theta)\alpha^{3/2}} \int \exp(-|\mathbf{r}' - \mathbf{r}_{12}|/\delta) K_\beta(\mathbf{r}') d\mathbf{r}',$$

$$\qquad (7)$$

$$\mathbf{r}_{12} = \mathbf{r}_1 - \mathbf{r}_2,$$

where $\delta = (\alpha/a)^{1/2}$ is the characteristic action radius of the exchange forces. The actual form of $K_\theta(\mathbf{r})$ depends substantially on the ratio of the characteristic lengths R_c and δ. In particular, at $R_c/\delta \ll 1$ we have

$$K_\theta(\mathbf{r}_{12}) = D(\beta) R_c^3 \delta\alpha^{-2} \exp(-|\mathbf{r}_{12}|/\delta) \sin^2\theta\cos^2\theta, \qquad (8)$$

i.e., the role of the correlation radius of the fluctuations $\delta\theta$ is assumed by the characteristic length δ. At $R_c/\delta \gg 1$ we have

$$K_\theta(\mathbf{r}_{12}) = D(\beta)\delta^4\alpha^{-2} \exp(-|\mathbf{r}_{12}|/R_c) \sin^2\theta\cos^2\theta, \qquad (9)$$

i.e., as expected, the correlation radius of the spin-flip fluctuations coincides in this case with the correlation radius of the anisotropy fluctuations. We note that according to (7)–(9) the statistical characteristics of the fluctuations $\delta\theta$ are anisotropic and depend on the orientation of the average magnetic moment of the system. The fluctuations $\delta\theta(\mathbf{r})$ can be regarded as small if $K_\theta(\mathbf{r}) \ll 1$ at distances of the order of the correlation radius. This yields[2)]

$$D(\beta)R_c^3/a^2\delta^3 \ll 1 \qquad (R_c/\delta \ll 1), \qquad (10)$$

$$D(\beta)/a^2 \ll 1 \qquad (R_c/\delta \gg 1). \qquad (11)$$

Thus, in the first case large fluctuations of the anisotropy with a correlation radius small compared with δ are admissible; such a situation can be observed for example for a random distribution of strongly anisotropic impurities. The second case can be likened to extended defects with small variations of the anisotropy constants.

We transform now the expression for the density of the thermodynamic potential of the system. Recognizing that in accord with (7)–(9) the $\delta\theta(\mathbf{r})$ random field is statistically homogeneous, and has the property that the correlations vanish at infinitely remote points, and using the ergodic theorem[8]

$$\lim_{V \to \infty} V^{-1} \int \dots dV = \langle \dots \rangle,$$

we obtain

$$\Omega_{eff} = -\tfrac{1}{2}\beta\cos^2\theta + \tfrac{1}{2}\alpha\langle(\nabla\theta)^2\rangle$$

$$+ \tfrac{1}{2}\beta\cos 2\theta\langle\delta\theta^2\rangle + \sin\theta\cos\theta\langle\delta\beta\delta\theta\rangle. \qquad (12)$$

If we confine ourselves in (12) to the first term, which takes into account only the average anisotropy energy of the system (this is precisely the assumption used to study transitions in disordered systems[1]), we obtain an approximation similar to the virtual-crystal approximation of the theory of

disordered materials.[9] The remaining terms in (12) take into account the contribution made to the energy by the fluctuations $\delta\theta$ inherent in the systems.

Substituting in (12) the values of $\delta\theta$ from (6), we obtain

$$\Omega_{eff} = -\frac{\beta}{2}\cos^2\theta - \frac{D(\beta)R_c^2}{2\alpha}\sin^2\theta\cos^2\theta\left[1 + \frac{R_c}{\delta}\right]^{-2}. \qquad (13)$$

If we neglect R_c/δ in the second term of (13), Ω_{eff} reduces to

$$\Omega_{eff} = K_1\sin^2\theta + K_2\sin^4\theta, \qquad (14)$$

$$K_1 = \beta/2 - D(\beta)R_c^2/2\alpha, \qquad (15a)$$

$$K_2 = D(\beta)R_c^2/2\alpha. \qquad (15b)$$

Thus, allowance for the fluctuations $\delta\theta$ due to the magnetic-anisotropy energy fluctuations leads to a renormalization of the thermodynamic potential of the system (when the latter is expressed in terms of $\bar\theta$) and to the appearance of effective anisotropy constants of higher order.

We note that despite the smallness of the fluctuation $\delta\theta$, their contribution to the effective thermodynamic potential can be appreciable. Thus, at

$$R_c/\delta \approx \beta/D(\beta)^{1/2}, \qquad \bar\theta \approx 0,$$

it follows from (5) that it is comparable with $\bar\beta$. The fluctuations $\delta\theta$ are in this case no less small, since relation (1) takes the form $\bar\beta/D(\beta)^{1/2} \ll 1$ and is well satisfied, e.g., in the spin-flip region, where the constant $\bar\beta$ is small.

Minimization of Ω_{eff} yields three equilibrium phases[1]: $\bar\theta = 0$, $\bar\theta = \pi/2$, and a canted one in which θ takes on intermediate values. For the thermodynamic potential (14), for example, we have in the canted phase the known relation $\sin^2\bar\theta = -K_1/2K_2$ (Ref. 1).

The regions of existence of the different phases are shown in Fig. 1, where we use for convenience the dimensionless parameters

$$z = D(\beta)R_c^4/\alpha^2, \qquad \bar y = R_c^2\beta/\alpha.$$

When $\bar y$ varies with temperature (this corresponds to a change of the average anisotropy constant), orientational transitions take place in the system at temperatures T_1 and T_2 defined by the relations
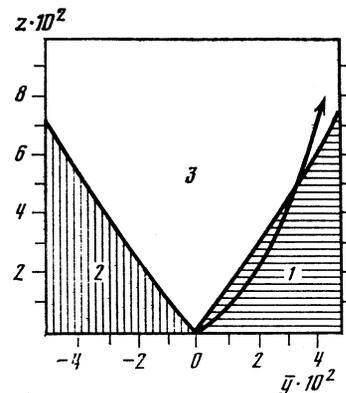
FIG. 1. Phase diagram. The arrow indicates the direction of motion for spin-flip transitions connected with the decrease of the relative exchange energy. 1—$\bar\theta = 0$, 2—$\bar\theta = \pi/2$, 3—canted phase.

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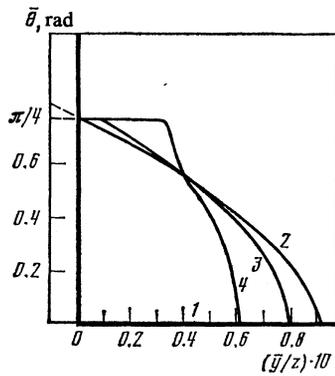
FIG. 2. Possible form of the dependence of $\bar{\theta}$ $(\bar{y})$ for different sections of the phase diagrams of Fig. 1: 1—ordered medium $z = 0$;2—$z = 10^{-3}$;3— $z = 10^{-2}$;4—$z = 0,1$.

$$z = |\bar{y}(T_{1,2})| (1 + |\bar{y}(T_{1,2})|^{1/2}). \tag{16}$$

It can be shown that when conditions (10) and (11) are satisfied we have at these points the divergence of the averaged (over the fluctuations $\delta\beta$) generalized susceptibility in accord with a law characteristic of the Landau theory $\langle \chi \rangle \propto (T - T_{1,2})^{-1}$, i.e., in this sense these transitions are second-order phase transitions. The possible types of temperature dependence of the angle $\bar{\theta}$ in the spin-flip region are shown in Fig. 2. We note that with increasing z the size of the section $\theta \approx \pi/4$ at the center of the transition region increases, and the transition at the points T_1 and T_2 becomes more and more abrupt and acquires the character of a jump.

Interestingly, it is possible to have in the considered systems spin-flip transitions not connected with the change of the sign of the average anisotropy constant $\bar{\beta}$. Indeed, in accord with (15a), for example, the constant K_1 is determined by the contribution of two terms of opposite sign. When the exchange interaction is weakened (this can be observed for example, in rare-earth ion compounds, where the magnetic anisotropy energy exceeds the exchange when the temperature is lowered,[10,11] a negative contribution to K_1 may predominate. These transitions will be observed on the phase diagram of Fig. 1 when moving along the lines

$$\bar{y} = R_c^2 \bar{\beta}/\alpha, \quad z = D(\beta)R_c^4/\alpha^2 \tag{17}$$

as $\alpha \to 0$. The intersections of the curves of the family (17) with the phase-transition line (and consequently also the realization of spin-flip transitions due to weakening of the exchange) is possible, as can be readily shown, only at sufficiently large fluctuations of the magnetic-anisotropy energy, $D(\beta)/\beta^2 \gtrsim 1$. These transitions do not occur, and the system remains in the canted phase characterized by an angle $\bar{\theta}_{lim}$ whose value is determined by the statistical characteristics of the anisotropy fluctuations. At $\beta^2/D(\beta) \ll 1$, for example, $\bar{\theta}_{lim} \approx \pi/4$.

The conditions for the applicability of the described scheme of spin-flip transitions in systems with fluctuations of the magnetic-anisotropy energy can be determined on the basis of relations (10) and (11). It can be shown that the expressions (13) and (14) used by us for the density of the thermodynamic potential become rather crude approximations in the immediate vicinity of the line $\bar{y} = 0(\beta = 0)$, i.e., at the center of the transition region.

We present by way of example several estimates for disordered solid solutions on the basis of rare-earth orthoferrites, where $\alpha \approx 4 \times 10-7$ erg/cm^3 and the constant $\bar{\beta}$ can be represented in the form

$$\beta = K(T - T_r)/T_r = Kt, \tag{18}$$

where $K \approx 10^4$ erg/cm^3 (Ref. 1).

In the compounds considered there should be added to ω_{an} also the fourth-order anisotropy $\omega_{an}^{(4)} = k_2 \sin^4\theta$. Since the constant k_2 in orthoferrites is small ($k_2 \approx 10^3$ erg/cm^3), allowance for its fluctuations does not make substantial contributions to the thermodynamic potential (13) or (14), and in the latter we can confine ourselves to allowance for the average energy $\bar{k}_2 \sin^4\theta$. The characteristic effective radius of the exchange forces in compounds based on orthoferrites is large: $\delta \gtrsim 10^3 a$, where a is the lattice constant. We can therefore expect satisfaction of the condition $R_c/\delta \ll 1$ and, consequently, describe the systems by the thermodynamic potential (14) with effective constant

$$K_{2eff} = \bar{k}_2 + D(\beta)R_c^2/2\alpha = \bar{k}_2 + \Delta k_2.$$

In the case of substitutions in the rare-earth and iron sublattices (except for the cases when Fe^{3+} ions are replaced by the ions Co^{2+} and Mn^{3+}, which have an unusually large magnetic-anisotropy energy[1], the magnitude of the fluctuations $D(\beta)^{1/2}$ corresponds to 10^5–10^6 erg/cm^3 (Ref. 1), so that

$$\Delta k_2 \approx (10^1 - 10^3)(R_c/a)^2 \text{ erg/cm}^3$$

and even in fully disordered solid solutions ($R_c \approx a$) it can make a substantial contribution to K_{2eff}. When the correlation radius R_c is increased, say as a result of partial ordering of the solid solution, Δk_2 increases correspondingly.

Since Δk_2 is positive, one can expect at $\bar{k} < 0$ the resultant constant K_{2eff} to be positive, and consequently the second anisotropy constant in the disordered and in the corresponding ordered compound will have opposite signs. In the ordered compound the spin-flip takes place jump wise (first-order phase transition[1]). and in the disordered one we have two second-order phase transitions at the points T_1 and T_2 defined by the condition (16). The change of the character of the transition was experimentally observed[1] for a transition in the ab plane of YFeO$_3$ (constant $\bar{k}_2^{ab}$(YFeO$_3$) < 0, Ref. 11) when part of the Fe^{3+} ions was replaced by Mn^{3+} ions. There are also data[12] on the observation of a similar phenomenon in disordered solid solutions based on hematite $(Fe_{1-x}Cr_x)_2O_3$, $(Fe_{1-x}Al_x)_2O_3$.

At $k_2 > 0$ the fluctuations of the anisotropy increase the width of the transition. In fact, according to Ref. 1, in systems described by the thermodynamic potential (14), the temperatures of the start and of the end of the transition are determined by the relations $K_1(T_2) = 0, K_2(T_1) = -2K_2(T_1)$. From this we obtain for the ratio of the temperature intervals $\Delta T = T_2 - T_1$ in a disordered $(\Delta T(D(\beta))$ and ordered $(\Delta T(0))$ compound

$$\Delta T(D(\beta))/\Delta T(0) = 1 + \Delta k_2/k_2.$$

In accord with the estimates of Δk_2 above, the broadening of the transition region in a disordered compound can be appreciable. It is possible that this mechanism, alongside the peculiarities noted in Ref. 1 for the temperature dependence of the constant K_1, will lead to a large width of the transitions in a number of the solid solutions $YFe_{1-x}Cr_xO_3$.

In conclusion, we determine the applicability limits of the model used by us for the orientational transitions in the compounds considered. Substituting in (10) the equilibrium values of the angle $\bar\theta$, determined from the condition that (14) be a minimum, we obtain

$$D^{\prime\prime}(\beta)R_c{}^4/\alpha^2 \ll \beta;$$

whence, using the estimates given above for the parameters, we obtain

$$(10^{-2}-10)(R_c/a)^4\mathrm{erg/cm^3} \ll \beta.$$

This relation excludes temperatures in the immediate vicinity of the center of the spin-reorientation region $(\bar\theta = \pi/4, \beta = 0)$. Taking (18) into account we find that for disordered compounds based on orthoferrites our analysis is valid if the relative temperature t satisfies the condition $t \gg (10^{-6} - 10^{-3})(R_c/a)^4$. In the case of a completely disordered solid solution $(R_c \approx a)$ the region of the temperatures excluded by this condition is extremely small.

3. SPIN-FLIP IN COMPOUNDS WITH FLUCTUATIONS OF THE ANTISYMMETRIC EXCHANGE

Allowance for the interaction of the systems considered in the preceding section with an external magnetic field does not lead to additional, compared with (13), renormalizations of the thermodynamic potential. An entirely different picture is observed in the presence of fluctuations of the Zeeman energy of the system.

We consider by way of example spin-flip transitions in disordered compounds based on orthoferrites with fluctuations of the antisymmetric exchange. The latter is known[1] to lead to the onset of a transverse weak moment in the sublattice of the $3d$-ions. We consider for the sake of argument transitions in the ac plane, in a field directed along the c axis (Fig. 3). The density of the thermodynamic potential of the system can be represented in this case in the form[1]
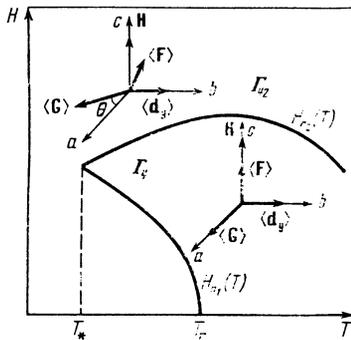
FIG. 3. Phase diagram of a compound with fluctuations of the antisymmetric exchange. The insets show the geometry of the arrangement of the vectors $\langle F \rangle$, $\langle G \rangle$, $\langle d_y \rangle$, $\langle H \rangle$ in the phases Γ_4, Γ_{42}, $(\chi_{\parallel} > \chi_{\perp})$.

$$\Omega = \lim_{V\to\infty} V^{-1}$$

$$\times \int \left\{ \frac{\alpha}{2}(\nabla G)^2 + \frac{I}{2}F^2 - d(r)[F\times G] - K_1G_x{}^2 - M_0FH \right\} dV, \tag{19}$$

where V, as above, is the volume of the system, and F and G are the mutually perpendicular ferro- and antiferromagnetism vectors.[1] The first and second terms in (19) take into account the exchange energy, and the third the antisymmetric Dzyaloshinskiĭ exchange.[6] In compounds base on $RFeO_3$ the random vector d is collinear with the b axis (Fig. 3).

We choose the correlation function $K_{d_y}(r-r')$ in the form (3) with the appropriate substitutions $D(\beta) \to D(d_y), R_c \to r_c$. The last two terms in (18) take into account the energies of the magnetic anisotropy and of the interaction with the external field; M_0 is the magnetization of the system at absolute saturation. For simplicity we shall disregard hereafter the fluctuations of the magnetic anisotropy. Eliminating F from (19) we obtain for $H \parallel c$ (Ref. 1)

$$\Omega = \lim V^{-1} \int \left\{ \frac{\alpha}{2}(\nabla G)^2 \right.$$

$$\left. - K_1G_x{}^2 - \frac{\chi_\perp h^2}{2}(1-G_z{}^2) - M_c hG_x \right\} dV, \tag{20}$$

$$\chi_\perp = 1/I, \quad h = M_0H_z, \quad M_c = d_y(r)/I.$$

The last term in (20) corresponds to fluctuations of the interation with the external field.

Introducing the angle $\theta(r)$ (Fig. 3) and eliminating, as in the preceding section, the fluctuations $\delta\theta$, we obtain

$$\Omega_{\rm eff} = \Omega_0 + \Delta\Omega, \tag{21}$$

where Ω_0 takes into account the average energy of the antisymmetric exchange and of the magnetic anisotropy[1]:

$$\Omega_0 = -K_1\bar{G}_x{}^2 - {}^1\!/_2\chi_\perp h^2(1-\bar{G}_z{}^2) - m_0h\bar{G}_x.$$

Here $\bar{G}_{x,z} = \langle G_{x,z} \rangle$ are the components of the average antiferromagnetism vector of the system, $m_0 = \langle d_y \rangle/I$, and $\Delta\Omega$ is a contribution due to the orientational fluctuations $\delta\theta$ [the conditions for its applicability can be established in analogy with (10) and (11)]

$$\Delta\Omega \approx -{}^1\!/_2\chi_\parallel h^2\bar{G}_z{}^2.$$

The quantity $\chi_\parallel$, determined by the fluctuations of the antisymmetric exchange, can be interpreted as the longitudinal susceptibility of the system:

$$\chi_\parallel = r_c{}^3 D(d_y)/\alpha I^2. \tag{22}$$

Minimization of $\Omega_{\rm eff}$ with respect to the angle $\bar\theta$ that specifies the orientation of the vector $\langle G \rangle$ yields two solutions (Fig. 3): the phase 4 (Ref. 1) in which

$$\sin\bar\theta = 0,$$

and the canted phase Γ_{42} (Ref. 1), in which

$$\cos\bar\theta = -\frac{\langle d_y \rangle h}{2I[K_1 + {}^1\!/_2(\chi_\perp - \chi_\parallel)h^2]}.$$

Investigation of the stability of the obtained solutions shows that the behavior of the considered systems depends

substantially on the sign $\Delta_\chi = \chi_\parallel - \chi_\perp$. At $\Delta_\chi > 0$, just as in ordered system, a transition $\Gamma_{42} \to \Gamma_4$ is observed at $K_1 < 0$ and $H \geqslant H_n$, with the threshold field determined by the relation[1]

$$H_n = -\frac{H_D}{2} + \left[\left(\frac{H_D}{2} \right)^2 - H_{an} H_{ex} \right]^{1/2}$$

with fields renormalized on account of the antisymmetric-exchange fluctuations

$$H_D = \langle d_y \rangle / M_0 \varepsilon, \qquad H_{an} = 2K_1 / M_0 \varepsilon,$$

where $\varepsilon \equiv |1 - r_c^2 D \langle d_y \rangle / \alpha I|$; the exchange field is of the usual form $H_{ex} = I/M_0$.

At $\Delta_\chi < 0$ two transitions are observed. The first is possible only at $K_1 < 0$ and corresponds to the usual $\Gamma_{42} \to \Gamma_4$ transition in an ordered compound. The threshold field is given by

$$H_{n_1} = \frac{H_D}{2} - \left[\left(\frac{H_D}{2} \right)^2 + H_{an} H_{ex} \right]^{1/2}.$$

Besides, at any sign of K_1 and when the condition $H_{an} > -H_D^2 / 4H_{ex}$ is satisfied, a transition $\Gamma_4 \to \Gamma_{42}$ is possible in a field

$$H_{n_2} = \frac{H_D}{2} + \left[\left(\frac{H_D}{2} \right)^2 + H_{an} H_{ex} \right]^{1/2}.$$

It is similar to the transition to the canted phase in ferrimagnets (for example, at $K_1 = 0$ we have the relations $\langle F_a \rangle = 0$, $\langle F_c \rangle = \chi_\parallel h$) that are usual for ferrimagnets), and is due to the presence in the system of a certain fraction of moments directed, owing to fluctuations of the antisymmetric exchange, counter to the weak moment. At $\mathbf{H} \parallel \langle \mathbf{F} \rangle$ deflections of these moments are possible toward the field that causes, owing to the exchange coupling spin flip in the neighboring regions and in the system on the average. Favoring the considered transition are large values of the correlation radius r_c. Indeed, the necessary condition ($\Delta_\chi < 0$) for the transition is transformed in the presence of the relation $\alpha \approx I a^2$ into

$$r_c^2 / a^2 \gtrsim I^2 / D \langle d_y \rangle,$$

inasmuch as in compounds based on orthoferrites we have $I/D\langle d_y \rangle \gg 1$ and the transition considered is possible only for large-scale ($r_c \gg a$) fluctuations of $d_y(\mathbf{r})$.

Figure 3 shows the phase diagram of the considered compounds under the assumption that $K_1 > 0$ at high temperatures and reverses sign at the temperature T_r. We note the possibility of the existence of a critical temperature T_* such that at $T < T_*$ no spin flip takes place in an external field, i.e., the critical transition field does not exist at all. The temperature T_* is determined by the condition $H_{n1} = H_{n2}$ or, equivalently, by the relation

$$m_0^2(T_*) / 4\Delta\chi = |2K_1(T_*)|.$$

We note in conclusion that an alternative to the considered transition in the field H_{n_2} is rotation of the vector $\mathbf{G}$ in the plane $\mathbf{ac}$; this rotation is known[1] to decrease the weak magnetic moment. In the presence of moments directed counter to the field as a result of fluctuations of $d_y(\mathbf{r})$, this process may turn out to be energetically preferred. An expression for the corresponding threshold field can be easily obtained in analogy with the preceding one. Thus, spin-flip in the $\mathbf{ac}$ plane in disordered compounds is possible only in the presence of large anisotropy, which holds the vector $\mathbf{G}$ in this plane. In the opposite case, transitions can be realized with spatial rotation of the vector $\mathbf{G}$.

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[1] Allowance for the fluctuations of α does not lead to qualitative peculiarities of the spin-flip transitions in view of the isotropic character of the exchange.

[2] It can be shown that the conditions under which the term $\delta\beta\delta\theta$ in Eq. (5) can be neglected take the form (10), (11) with inessential additional numerical factors in the left-hand sides.

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