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Nature of localized magnetic moments in band antiferromagnets

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It is shown that a nonmagnetic interaction of an impuity atom with a spin-density wave of band electrons of an antiferromagnet leads to formation of an electron state, localized on the impurity, with an uncompensated spin. The phenomenon takes place because the interaction with the impurity violates the equivalence of the electronic spin subbands and is analogous in many respects of excitonic ferromagnetism. The localized magnetic field is calculated as a function of the position of the Fermi level. The results of the model can be easily applied to the case of exchange interaction between a magnetic impurity and a charge-density wave, and can describe the screening of the bare magnetic moment of the impurity by the spins of the band electrons (in analogy with the Kondo effect). The proposed theory explains the nature of the magnetic moments of certain impurities of other lattice defects. The theory is compared with experimental data obtained for chromium alloys.

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1. We propose in this paper a model for the production of a localized magnetic moment in systems with electron–hole pairing.[1,2] In these systems, as a result of the special properties of the electron spectrum (the presence of congruent sections of the Fermi surfaces of the electrons and holes), even an arbitrarily weak electron–hole interaction leads, with decrease of temperature, to substantial changes in the electron subsystem. Below a certain temperature T_c a long-range order is produced in the system and is characterized by a parameter Δ. The value of Δ is proportional to the density of the condensate of the electron–hole pairs.

Depending on the nature of the electron–hole interaction, different types of spin structure of the ordered state are possible.[2] In the case of singlet pairing, the order parameter Δ_s is independent of the spin and corresponds to formation of a charge density wave (CDW) in the system. Pairing of electrons in the triplet state leads to formation of a spin-density wave (SDW). In this case the spin structure of the order parameter takes the form $\Delta_t=(\Delta\cdot\sigma_{\alpha\beta})$, where $\sigma_{\alpha\beta}$ are 2×2 Pauli matrices. It is known[3,4] that the formation of the SDW is reflected in the antiferromagnetic behavior of chromium.

Introduction of an impurity atom in a system with SDW leads to a local redistribution of the spin of the band electrons such that in the vicinity of the impurity there appears a nonzero magnetic moment. When analysing the influence of the impurity on the electron subsystem of the metal, a distinction must be made between two types of interaction of the electrons with the impurity atom.

If the impurity does not have a magnetic moment, then its interaction with the band electrons is described by the usual potential scattering $V(\mathbf{r})\Psi_\alpha+\Psi_\alpha(\Psi_\alpha^+$ is the operator of production of a particle with spin α).

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If an uncompensated electron spin S is present on the inner electron shell of the impurity atom, then in addition to the potential scattering V there appears in the interaction of the electrons with the impurity an exchange interaction $J(\mathbf{r})S\hat{\sigma}_{\alpha\beta}\Psi_\alpha^+\Psi_\beta$.

It will be shown in Sec. 2 below that in a system with SDW the potential nonmagnetic scattering V of the electrons by the impurity leads to formation of a localized bound state for the electrons with a preferred spin. A similar bound state for the spin can arise for a magnetic impurity in the case of exchange interaction $J(\mathbf{r})S\hat{\sigma}$ with the charge-density wave. Depending on the sign $J(\mathbf{r})$, the bare atomic angular momentum S will either be cancelled out or be supplemented by the spin of the electron of the bound state. In this sense, this phenomenon is similar to the Kondo effect.[5]

In a system with strong correlation of the electrons and hole, simultaneous allowance for the potential and exchange scattering of the electrons by the impurity can lead to a screening of the magnetic moment even above the temperature at which a spin (charge) density wave is produced in the pure system. In fact, the exchange interaction of the band electrons with the impurity polarizes the electrons by spin, which corresponds formally to the appearance of a local triplet parameter $\Delta_t \sim J(\mathbf{r})S\langle\Psi_\alpha^+\Psi_\alpha\rangle$. Next, allowance for the potential scattering V by the same impurity atom, but now surrounded by the local spin density wave $\Delta_t \neq 0$ leads to the appearance of an excess spin of the electrons in the vicinity of the impurity.

It should be noted that the effects described above are similar to the phenomenon of excitonic ferromagnetism,[6] where the ferromagnetic ordering of the spins of the band electrons proceeds to the extent that the SDW ($\Delta_t \neq 0$) coexists with the charge-density wave ($\Delta_s \neq 0$) of the electrons. In the case considered by us the local source of one of these waves is the exchange or potential interaction of the electrons with the impurity atom.

2. We consider now the simplest model of an isotropic semimetal with identical electron and hole masses. As a result of the electron–hole interaction, such a system undergoes with decreasing temperature a phase transition into a state of an excitonic dielectric.[1] As already noted, this can produce in the system a charge density wave or a spin density wave. Hereafter we shall consider in the calculations the second case (SDW).

In the electron representation, for both bands the system is described below the phase-transition temperature T_N by the effective Hamiltonian

$$\mathcal{H}_0 = \sum_\sigma \int \left\{ \Psi_{1\sigma}^+(\mathbf{r}) \left(-\frac{\nabla^2}{2m}\right) \Psi_{1\sigma}(\mathbf{r}) - \Psi_{2\sigma}^+(\mathbf{r}) \left(-\frac{\nabla^2}{2m}\right) \Psi_{2\sigma}(\mathbf{r}) \right\} d\mathbf{r} + \sum_{\sigma\sigma'} \int \Delta^{\sigma\sigma'}(\mathbf{r}) \Psi_{1\sigma}^+(\mathbf{r}) \Psi_{2\sigma'}(\mathbf{r}) d\mathbf{r} + \text{c.c.}, \tag{1}$$

$\Psi_{1,2\sigma}^+(\mathbf{r})$ is the operator of production of a particle from band 1 or 2 with spin σ at the point $\mathbf{r}$.

In the Hamiltonian (1), the electron–hole interaction is written in the form of an interband Hartree-Fock potential $\Delta^{\sigma\sigma'}(\mathbf{r})$. As was done by Kopaev and Rusinov,[7] we shall assume that $\Delta(\mathbf{r})$ does not depend on the coordinate.

The spin structure of the potential $\Delta^{\sigma\sigma'}$ is chosen in accordance with the triplet-pairing condition

$$\Delta^{\sigma\sigma'} = \Delta\hat{\sigma}_z, \quad \hat{\sigma}_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$$

is a Pauli matrix.

To describe the interaction of the electrons with the impurity atom, we confine ourselves to the case of potential scattering by a pointlike potential $V(\mathbf{r}) = V\delta(\mathbf{r})$. To simplify the calculations it suffices to retain in this interaction only the terms corresponding to interband impurity scattering. It can be shown that simultaneous account of the intraband scattering does not change the qualitative results of the model (see Sec. 3). In this approximation, the interaction with the impurity is described by the Hamiltonian

$$\mathcal{H}_{int} = -\sum_\sigma \int V_{12}(\mathbf{r}) \Psi_{1\sigma}^+(\mathbf{r}) \Psi_{2\sigma}(\mathbf{r}) d\mathbf{r} + \text{c.c.} \tag{2}$$

In accordance with the Hamiltonian (1) and (2), the equations of motion for the temperature Green's functions

$$G_{\alpha\beta}^{\sigma\sigma'}(x, x') = \langle T\tilde{\Psi}_{\alpha\sigma}(x)\tilde{\Psi}_{\beta\sigma'}^+(x')\rangle, \quad x = (\mathbf{r}, t)$$

take the form

$$\begin{aligned} (i\omega + \hat{\varepsilon}(\mathbf{p}) + \mu) G_{11}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}') + \Delta\hat{\sigma}_z G_{21}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}') &= \delta(\mathbf{r} - \mathbf{r}') + V_{12}(\mathbf{r}) G_{21}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}'), \\ (i\omega - \hat{\varepsilon}(\mathbf{p}) + \mu) G_{21}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}') + \Delta^*\hat{\sigma}_z G_{11}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}') &= V_{21}(\mathbf{r}) G_{11}^{\sigma\sigma'}(\mathbf{r}, \mathbf{r}'). \end{aligned} \tag{3}$$

In momentum space, the complete Green's function $G_{\alpha\beta}^{\sigma\sigma'}(\mathbf{p}, \mathbf{p}')$ can be expressed in terms of the free Green's function $G_{\alpha\beta}^{(0)\sigma\sigma'}(\mathbf{p})$ (excitonic dielectric without impurity[1]) and the complete vertex for the scattering of electrons by an impurity center:

$$G_{\alpha\beta}^{\sigma\sigma'}(\mathbf{p}, \mathbf{p}') = G_{\alpha\beta}^{(0)\sigma\sigma'}(\mathbf{p})\delta_{\mathbf{p}\mathbf{p}'} + \sum_{\tau,\tau',\sigma''} G_{\alpha\tau}^{(0)\sigma\sigma''}(\mathbf{p}) \Gamma_{\tau\tau'}^{\sigma''\sigma''} G_{\tau'\beta}^{(0)\sigma''\sigma'}(\mathbf{p}'), \tag{4}$$

$$G_{\alpha\beta}^{(0)\sigma\sigma'} = \frac{1}{\omega^2 + \varepsilon^2 + \Delta^2} \begin{pmatrix} -i\omega - \varepsilon & \Delta\hat{\sigma}_z \\ \Delta^*\hat{\sigma}_z & -i\omega + \varepsilon \end{pmatrix}.$$

In the assumed approximation of a pointlike impurity potential we can obtain for the matrix of the complete vertex $\Gamma_{\alpha\beta}^{\sigma\sigma}$ the expression

$$\Gamma_{\tau\tau'}^{\sigma\sigma} = \frac{1}{D_0} \begin{bmatrix} -\mathcal{V}^2 i\omega_- & \mathcal{V} d_0 - \mathcal{V}^2(\Delta\hat{\sigma}_z) \\ \mathcal{V} d_0 - \mathcal{V}^2(\Delta\hat{\sigma}_z) & -\mathcal{V}^2 i\omega_- \end{bmatrix} \tag{5}$$

$$\mathcal{V} = \pi N(\varepsilon_F) V, \quad \omega_- = \omega - i\mu, \quad d_0 = (\omega_-^2 + \Delta^2)^{1/2}.$$

Here $N(\varepsilon_F)$ is the state density in the semimetallic phase at the Fermi level,

$$D_0 = (1 + \mathcal{V}^2)(\omega_-^2 + \Delta^2)^{1/2} - 2\mathcal{V}(\Delta\hat{\sigma}_z). \tag{5a}$$

From expression (5a) for the denominator D_0 we see that in the case when $\hat{\sigma}_z = +1$ the vertex part has two imaginary poles

$$\omega_0 = \pm\Delta \frac{1 - \mathcal{V}^2}{1 + \mathcal{V}^2}. \tag{6}$$

This means that two local impurity levels $|\omega_0| < |\Delta|$, appear in the forbidden band of the excitonic dielectric, and both levels are allowed only for electrons with definite and identical spin projections ($\sigma = +1$).

It is useful to note for future reference that inasmuch

as the total number of states in the system should not change, it follows that the levels $\pm\omega_0$ must be regarded as spatially separated from the valence band ($\omega = -\omega_0$) and from the conduction band ($\omega = +\omega_0$).

The choice of the corresponding spin projection is determined by the sign of the product of the interband element of the impurity scattering V_{12} by the quantity Δ. If the impurity lands in a site of the host lattice, then the magnitude and the sign of the matrix element are determined by the Bloch functions of the electrons and do not change from site to site. As to the sign of the parameter Δ, it can oscillate. For example, in the characteristic case of chromium the SDW has a period twice as large as that of the host lattice.[3] The neighboring sites of the host lattice are in antinodes of the SDW with opposite signs for Δ. Therefore the direction of the spin localized on an impurity atom depends on the site on which this atom lands.

If the concentration of the impurities in this system is small (so that their interaction can be neglected), then it is reasonable to assume that on the average a number of impurity atoms with positive localized moment is equal to the number of atoms with negative moment, i.e., the average magnetic moment of the system is zero.

We turn now to the value of the localized spin moment. With the aid of the expressions for the total Green's function (4) and the total vertex (5) we can calculate the increments for the total number of particles of one spin subband:

$$\delta N^{\sigma} = T\sum_{n}\int [\delta G_{11}^{\sigma\sigma}(\mathbf{p},\mathbf{p}) - \delta G_{22}^{\sigma\sigma}(\mathbf{p},\mathbf{p})]\frac{d^3p}{(2\pi)^3}. \tag{7}$$

Then the total change of the spin density δS and of the particle-number density δN in the system is determined by

$$\delta S = {}^1/_2[\delta N^{\sigma} - \delta N^{-\sigma}], \quad \delta N = \delta N^{\sigma} + \delta N^{-\sigma}. \tag{8}$$

Substituting the expression for $\delta G_{\alpha\beta}^{\sigma\sigma}(\mathbf{p},\mathbf{p})$ from (4) in (7) and (8) and leaving out the cumbersome calculations, we obtain

$$\delta S = -\frac{1}{2}[n(-\mu+\omega_0)+n(-\mu-\omega_0)] + \frac{2\tilde{\Gamma}\Delta}{\pi(1+\tilde{\Gamma}^2)}\int_{\Delta}^{\infty}\frac{\omega[n(-\mu+\omega)+n(-\mu-\omega)]}{(\omega^2-\Delta^2)^{1/2}(\omega^2-\omega_0^2)}d\omega, \tag{9}$$

$$\delta N = -[n(-\mu+\omega_0)+n(-\mu-\omega_0)], \quad n(\omega) = [1+e^{\omega/2T}]^{-1}. \tag{10}$$

In expression (9) for δS the first term corresponds to the spin filling the corresponding local level in the forbidden band, while the second term, represented by the integral, describes the polarization of the spins of the band electrons. The latter can be clearly illustrated with the aid of the following simple reasoning. Near the impurity there is a local change of the spectrum for each spin subband. For the up-spins ($\sigma = +1$) the effective gap in the spectrum decreases in proportion to $|\Delta - V_{12}|$, and for the spins with projection down ($\sigma = -1$) it increases in proportion to $|\Delta + V_{12}|$. Therefore the quasiparticles with up-spins accumulate near the impurity, and those with down-spins, on the contrary, are crowded out of this region. An uncompensated electron spin thus appears in the region of the impurity.

At $T = 0$ the integral in (9) can be exactly calculated and it turns out to be

$$\frac{1}{2} - \frac{1}{\pi}\operatorname{arctg}\left(\frac{\mu^2-\Delta^2}{\Delta^2-\omega_0^2}\right)^{1/2} f(\mu), \tag{11}$$

$$f(\mu) = \begin{cases} -1, & \mu > \Delta \\ 0, & -\Delta < \mu < \Delta. \\ 1, & \mu < -\Delta \end{cases} \tag{12}$$

It is convenient to represent the dependence of δS on the position of the Fermi level μ graphically (see the figure). It is seen from the figure that the magnetic moment appears in the system only when the Fermi level lies below the local level $\omega = -\omega_0$ (or above the level $\omega = +\omega_0$). The reason for the physical similarity of the picture is that the level $\omega = -\omega_0$ is split off, as already mentioned, from the states of the valence band. Therefore the system as a whole is nonmagnetic when all the valence-band states (including also the bound state separated from it) are filled with electrons. This situation is fully equivalent to that in an excitonic ferromagnet, the ferromagnetic moment vanishes when both spin subbands become filled.

Analyzing the expression for the total change of the number of particles in the system at $T = 0$, we can show that if there is no extra electron in the system, then the Fermi level lands in the interval $-\Delta < \mu < -\omega_0$, and in accordance with (9) (see the figure) a magnetic moment appears, the gist of which is the polarization of the spins of the hole-band electrons.

In the case of incomplete congruence of the Fermi surfaces of the electrons and holes, partial dielectrization of the carriers takes place. The remaining free carriers can determine to a considerable degree the position of the Fermi level. Then the value of δS will depend on the magnitude and sign of the doping, in accordance with (9).

Using expressions (4) and (5) for the total off-diagonal Green's function $G_{12}^{\sigma\sigma}(\mathbf{p},\mathbf{p}')$, we can also determine the corrections to the parameter. On the average over the volume Ω of the system we obtain

$$\frac{\delta\Delta^{\sigma\sigma}}{g_i} = \frac{1}{\Omega}\left\{\frac{1+\hat{\sigma}_z}{2}\frac{\omega_0}{4\Delta}\left(\operatorname{th}\frac{\mu-\omega_0}{2T} - \operatorname{th}\frac{\mu+\omega_0}{2T}\right) - \frac{\tilde{\Gamma}}{\pi(1+\tilde{\Gamma}^2)}\left[\frac{\omega_0^2}{2}\int_{\Delta}^{\infty}\frac{d\omega}{(\omega^2-\Delta^2)^{1/2}(\omega^2-\omega_0^2)}\left(\operatorname{th}\frac{\mu-\omega}{2T} - \operatorname{th}\frac{\mu+\omega}{2T}\right)\right.\right.$$
$$\left.\left. + \int_{-\varepsilon_F}^{\varepsilon_F}\frac{d\varepsilon}{(\varepsilon^2+\Delta^2)^{1/2}}\left(\operatorname{th}\frac{\mu+(\varepsilon^2+\Delta^2)^{1/2}}{2T} - \operatorname{th}\frac{\mu-(\varepsilon^2+\Delta^2)^{1/2}}{2T}\right)\right]\right\}. \tag{13}$$

The first term in (13) contains the spin dependence and reflects the fact that the introduction of the nonmagnetic

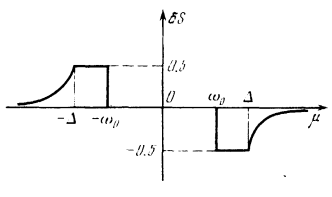

FIG. 1.

impurity into a system with a SDW violates locally the equivalence of the different spin subbands.

3. The results above pertain to the simplest case of nonmagnetic impurity in a system with SDW ($\Delta_t = \Delta\sigma_z$), but can be readily generalized to include other cases. In particular, these results, apart from certain signs, can be completely transferred to the case of a magnetic impurity ($J \neq 0$) in a system with a CDW ($\Delta_s \neq 0$). In this case, as indicated in Sec. 1, the bare magnetic moment of the impurity changes substantially to the extent of its screening by the spins of the band electrons. We shall henceforth neglect everywhere in the calculations the intraband scattering by the impurity V_{11} compared with the interband scattering V_{12}. The opposite limiting case $V_{12} \equiv 0$ is treated by Kopaev and Rusinov. They have shown that for an attracting (repelling) center in the forbidden band of an excitonic dielectric there is produced one local donor (acceptor) level that is degenerate in spin regardless of the spin structure of the initial state.

If we introduce interband scattering successively into this system with SDW, then we can show that such a localized level splits into two spin sublevels $\omega_{1,2}$ satisfying the equation

$$(1+V_{12}^2-V_{11}^2)(\Delta^2-\omega^2)^{1/2}+2[V_{11}\omega-V_{12}(\Delta\hat{\sigma}_z)]=0. \quad (14)$$

The main parameter in our problem is the interband scattering of the electrons V_{12}, but the latter can be due not only to a scattering by the impurity atom. For example, this scattering can be the result of other arbitrary lattice defects: vacancies, dislocations, or phase separation surfaces. This fact extends greatly the possibilities of the described model of localized magnetic states, and may possibly explain the nature of magnetic moments at such defects. Usually the expression of the nature of formation of localized magnetic moments at impurity atoms reduces to the most widely used Anderson model.[8] In this model the spin moment appears to the extent that the exchange Coulomb interaction U_{cor} of the electrons on the inner electron shell of the atom is strong enough. On account of the covalent interaction V of the electrons of the inner shell of the atom with the band electrons, this spin state can break down under the condition

$$U_{cor}/N(\varepsilon_F)V<1. \quad (15)$$

Anderson's theory explains successfully the nature of formation of certain magnetic moments, but is subject to definite limitations. For example, Anderson's theory cannot be used for solvent metals that exhibit noticeable superconducting properties, or for metals that undergo a metal–insulator transition [$N(\varepsilon_F)$ is large].

The fact that the localized moments in Cr–Fe alloys cannot be described by Anderson's simple model was recently noted by Angel *et al.*[9] As a result of an antiferromagnetic transition in chromium, the state density near the Fermi level $N(\varepsilon_F)$ increases substantially and hinders the localization in the Anderson model. Angel *et al.*[9] note that these difficulties can be overcome by modifying slightly the Anderson model for a two-band model of chromium.

The model of localized magnetic states proposed here differs significantly from the Anderson localization.[8,9] Whereas Anderson's theory is applicable for impurities of transition metals, in our model the nature of the impurity (defect) does not play a decisive role. It should be noted in this connection that the approach described in the present paper may possibly explain the measurements of the magnetic and electric properties of alloys of chromium with nontransition metals such as Al, Ge, Ga, Si, and Au.[10–14]

The physical difference between the model proposed here and the Anderson model is that in the latter the spin localization takes place on the inner shell of the impurity atoms, whereas in our model the local changes of the moment of the impurity are due to redistribution of the spin states of the band electrons surrounding the impurity, and to the formation of a bound state for the spin. The radius of this bound state depends on the depth of the local level ω_0 (6) and can be quite large.

In one of the experimental studies of the localized magnetic moments in a Cr–Fe alloy[15] with the aid of the Mossbauer spectra, the effective magnetic field H_{eff} at the ^{57}Fe nucleus was measured below the Neel point T_N. It is noted that H_{eff} has the same value as for the free atom. This allows us to conclude that the additional moment at the Fe impurity, which appears below T_N in this alloy, is not connected with localization of the excess spin on the inner shell of the Fe atom. Makarov *et al.*[15] point to the possibility of formation, around the Fe impurity, of a localized cloud of band electrons with polarized spin, a fact that has a direct bearing on the localized-moment model proposed here. Similar phenomena are apparently observed in the antiferromagnetic alloys Cr–Co (Refs. 16–19) and Cr–Mn (Refs. 20 and 21).

The presence of a long-range order (SDW or CDW) in the described model makes it possible to obtain in a self-consistent manner both a renormalization of the localized magnetic moment of the impurities and a correlation of these moments with one another. For example, in chromium, to the extent that the Fermi surfaces are not completely congruent, unpaired electrons are present and can carry the RKKY interaction[22] between the localized moments. Depending on the number of these free electrons, which is determined by the temperature and concentration of the impurity, the RKKY correlation of the moments will vary substantially.

In practically all the experiments[9–21], the susceptibility of the alloys is given by

$$\chi=\chi_0+C/(T-\Theta), \quad (16)$$

where χ_0 is the antiferromagnetic susceptibility of the pure chromium, while the second term in (16) corresponds to the susceptibility of the localized moments. The temperature Θ in this term characterizes the value of the magnetic interaction of the localized moments. As indicated, for example, by Hedgcock *et al.*, for the alloy Cr–Fe (% Fe ~ 1,2) the value of Θ changes dras-

tically from -80 ± 40 K for $T>T_N$ to -10 ± 5 K for $T<T_N$. With changing impurity concentration in the same alloy (at $T<T_N$) the value of Θ can change from -60 ± 10 K at $n_{Fe}\sim0.5\%$ to $+6\pm2$ K at $n_{Fe}\sim3.5\%$.[15]

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Ferrimagnets with Ising ions. Magnetic properties of holmium-yttrium iron garnets in strong fields at helium temperatures

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The magnetization of single-crystal holmium-yttrium iron garnets $Ho_xY_{3-x}Fe_5O_{12}$ ($0\leq x\leq3$) was measured in fields up to 300 kOe at 4.2 K. It was observed that phase transitions into field-induced noncollinear phases in compositions with $x<1.65$ are of first order H–x phase diagrams of holmium-yttrium iron garnets are constructed for field orientations along the crystallographic axes ⟨111⟩, ⟨110⟩, and ⟨100⟩. It is shown that the experimental data (the magnetization curves along different directions, the dependences of the transition fields on the holmium concentration and on the orientation of the magnetic field, etc.) of ferrites with $x<1.65$ can be explained by means of a model that considers the Ho^{3+} ion in the garnet structure in an extremely anisotropic Ising approximation. Possible causes of deviation from this model in iron garnets with large holmium contents are discussed.

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1. INTRODUCTION

Numerous recent investigations (see, e.g., the review[1]) have shown that noncollinear magnetic structures can be produced in ferrimagnets in a certain field interval. The formation of such structures is due to the competition between the negative exchange interaction that tends to produce an antiparallel alignment of the magnetic moments of the sublattices of the ferrimagnet, and the Zeeman interaction of the magnetic moment of the sublattices with the external field, which tends to orient them parallel to each other.

The field-induced noncollinear magnetic structures were investigated in greatest detail theoretically and experimentally in ferrimagnets in which the magnetic anisotropy is either nonexistent (the isotropic case)[2] or is small compared with the exchange interaction between the sublattices.[3] Much less investigated were the noncollinear structures in strongly anisotropic ferrimagnets. Anomalies of the magnetization of holmium–yttrium iron garnets (HYIG) $HO_xY_{3-x}Fe_5O_{12}$($x\lesssim0.4$) were observed in Ref. 4 in fields close to the exchange field (~10^5 Oe). It was shown theoretically[5] that these anomalies can be interpreted as a manifestation of the realignment of the magnetic structure of the ferrimagnet in the field. The rare-earth holmium ions were treated in this case in the extremely anisotropic Ising approximation.

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