Dynamics of amorphous magnets with strong random anisotropy

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The dynamics of amorphous magnets with strong random anisotropy is considered in the ferromagnetic-correlation region $\tau = (T - T_1)/T_1 \gg \tau_c$, where $T_1 = (c/3) \int J(r) d^3r$ and $\tau_c \approx (\varkappa^3/2c)^2$ [$\varkappa^{-1}$ is the $J(r)$ interaction radius and c is the density of the magnetic ions]. It is assumed that $c\varkappa^{-3} \gg 1$. It is shown by the dynamic-functional method that in the principal approximation the dynamic susceptibility takes the form $G^{-1}(\mathbf{k},\omega) \propto [(i\omega/\Gamma_0) + \tau + k^2/\varkappa^2]$, which is typical of purely dissipative dynamics. The corrections to $G(\mathbf{k},\omega)$ necessitated by application of an external magnetic field are calculated. There is no ferromagnetic resonance.

1. INTRODUCTION

From among the many problems related to spin glass, interest attaches to the study of amorphous magnetic alloys. The configurational disorder in these systems leads to the appearance of single-ion anisotropy with random orientation of the axes. The Hamiltonian usually employed in such cases was proposed by Harris, Pliske, and Zuckermann,[1] and is of the form

$$H[S] = -\frac{1}{2}\sum_{ij} J(r_{ij})\mathbf{S}_i\mathbf{S}_j - D\sum_i (\mathbf{S}_i\boldsymbol{\xi}_i)^2 - \mathbf{h}\sum_i \mathbf{S}_i. \quad (1.1)$$

Here $\mathbf{S}_i$ is the Heisenberg spin in the site i; ξ_i is a random unit vector that specifies the anisotropy direction, and $\mathbf{h} = g\mu_B\vec{\mathscr{H}}$ is the reduced magnetic field. An arbitrarily weak anisotropy ($D \ll J$) is known[2] to destroy the ferromagnetic state of such a system (at $h = 0$). The opposite case ($D \gg J$) was investigated experimentally in Refs. 3–7 and theoretically in Ref. 8. We have suggested in Ref. 8 that the radius over which the $J(r)$ interaction falls off is large compared with the distance a between neighboring spins. It was shown that the spin-correlation radius in the paramagnetic phase increases with decrease of temperature in the region $\tau \gtrsim \tau_c \approx (\varkappa^3/c)^2$, where $\tau = (T - T_1)/T_1$, $T_1 = (c/3)\int J(r) d^3r$, and c is the density of the magnetic ions; as a rule, $c \approx 0.5$. If $D \gg J$ each spin is directed along its easy magnetization axis and the microscopic variables become the Ising variables. On a macroscopic scale, however, the system acquires a soft mode constituting a molecular field that acts on the spin in a given site. This mode, which has vector properties, determines the large-scale static spin correlations. In the region $\tau \gtrsim \tau_c$ these correlations are of the same form as for a pure Heisenberg ferromagnet.[8] [Note that the maximum ferromagnetic-correlation radius $L_c \propto \varkappa^{-1}(\varkappa^3/c)^{-1}$ is substantially larger than the interaction radius.] As the temperature is reduced further, the growth of the correlation radius saturates and the Ising spin-glass regime sets in.[8]

We consider here the dynamics of amorphous magnets with strong random anisotropy, $D \gg J$, in the ferromagnetic-correlation region. Over long times and large scales this behavior is determined also by the above-mentioned soft mode. Using the dynamic-functional method proposed in Ref. 9, we obtain expressions for the susceptibility $\chi(\mathbf{k},\omega)$, which has in the principal approximation the form $\chi(\mathbf{k},\omega) \propto (i\omega + \tau + k^2)^{-1}$ typical of purely dissipative dynamics. The dynamic behavior of amorphous magnets in the feromagnetic-correlation region $\tau \gtrsim \tau_c$ is thus substantially different from that of Heisenberg ferromagnets.

We consider also the change of the form of $\chi(\mathbf{k},\omega)$ in the presence of an external magnetic field and show that the paramagnetic resonance typical of ferromagnets is absent in this case. The reason is the strong anisotropy which prevents spin precession in the magnetic field.

2. DERIVATION OF DYNAMIC FUNCTIONAL

At $D \gg J$ it can be assumed that each spin is directed along its easy-magnetization axis: $\mathbf{S} = \sigma\boldsymbol{\xi}_i$, where $\sigma_i = \pm 1$. We then obtain from (1.1) the Hamiltonian of the Ising variables σ_i:

$$H[\sigma] = -\frac{1}{2}\sum_{ij} J(r_{ij})\boldsymbol{\xi}_i\boldsymbol{\xi}_j\sigma_i\sigma_j - \mathbf{h}\sum_i \boldsymbol{\xi}_i\sigma_i. \quad (2.1)$$

The vectors $\boldsymbol{\xi}_i$ are assumed here to be uncorrelated:

$$\overline{\xi_i^\alpha \xi_j^\beta} = {}^1/_3\delta^{\alpha\beta}\delta_{ij} \quad (2.2)$$

(the superior bar denotes configurational averaging). The interaction radius $\varkappa^{-1}$ is defined as

$$\varkappa^{-2} = \frac{1}{6J_0}\int r^2 J(r)\, d^3r, \qquad J_0 = \int J(r)\, d^3r. \quad (2.3)$$

We put $J_0 = 1$ hereafter and assume that $c\varkappa^{-3} \gg 1$.

We start from the dissipative dynamics of the Ising variables σ_i with the Hamiltonian (2.1). This dynamics is described by the Glauber equations for the dynamics of Ising spins σ_i:

$$\dot{P}\{\sigma_i\} = \Gamma_0 \sum_j \Big[-\exp\Big(-\frac{1}{T}\sum_k J_{jk}\boldsymbol{\xi}_j\boldsymbol{\xi}_k\sigma_j\sigma_k\Big) \times P\{\sigma_1, \ldots, \sigma_j, \ldots, \sigma_n\} + \exp\Big(\frac{1}{T}\sum_k J_{jk}\boldsymbol{\xi}_j\boldsymbol{\xi}_k\sigma_j\sigma_k\Big) P\{\sigma_1, \ldots, -\sigma_j, \ldots, \sigma_n\}\Big], \quad (2.4)$$

where $P\{\sigma_i\}$ is the probability of the spin configuration $\{\sigma_i\}$, and Γ_0 is the spin-flip frequency. To derive the dynamic functional it is more convenient, however, to regard the variable σ as continuous ($-\infty < \sigma < +\infty$) with Hamiltonian

$$\tilde{H}[\sigma] = -\frac{1}{2}\sum_{ij} J_{ij}\sigma_i\sigma_j\boldsymbol{\xi}_i\boldsymbol{\xi}_j - \mathbf{h}\sum_i \boldsymbol{\xi}_i\sigma_i + \sum_i V_0(\sigma_i),$$

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$$V_0(\sigma)=u(\sigma^2-1)^2. \tag{2.5}$$

The transition to Ising variables occurs as $u\to\infty$. We shall be interested hereafter in times much longer than Γ_0^{-1}.

The temporal evolution of the variable σ is described by the equation

$$\dot\sigma_i+\partial\bar H/\partial\sigma_i=\zeta_i(t),\quad \langle\zeta_i(t)\zeta_j(t')\rangle=2T\delta_{ij}\delta(t-t'). \tag{2.6}$$

We introduce now a dynamic functional $Z\{\mathbf{l}_i\}$ defined as

$$Z\{\mathbf{l}_i\}=\left\langle \exp\left(\int dt\sum_i \sigma_i\xi_i\mathbf{l}_i\right)\right\rangle$$
$$=\int D\zeta D\sigma\delta\left[\dot\sigma_i+\frac{\partial\bar H}{\partial\sigma_i}-\zeta_i(t)\right]$$
$$\times\exp\left\{-\int dt\sum_i\left(\frac{\zeta_i^2}{4T}+\sigma_i\xi_i\mathbf{l}_i\right)\right\}$$
$$\times\det\left(\frac{\partial}{\partial t}+\frac{1}{2}\frac{\partial^2\bar H}{\partial\sigma_i(t)\partial\sigma_j(t')}\right). \tag{2.7}$$

Representing the delta-function in (2.7) by an integral with respect to a pure imaginary variable $p(t)$ and integrating over the noise ζ, we obtain

$$Z\{\mathbf{l}_i\}=\int D\sigma Dp\exp\left\{\int dt\sum_i\left[p_i\left(\dot\sigma_i-\sum_j\xi_i\xi_jJ_{ij}\sigma_j\right.\right.\right.$$
$$\left.\left.\left.-\xi_i\mathbf{h}+\partial V_0/\partial\sigma_i\right)+Tp_i^2+\sigma_i\xi_i\mathbf{l}_i\right]\right\}. \tag{2.8}$$

As in all problems with large interaction radius, it is convenient to introduce here continuous fields $\mathbf{f}(\mathbf{r})$ and $\mathbf{g}(\mathbf{r})$:

$$Z\{\mathbf{l}_i\}=\int D\sigma Dp D\mathbf{f} D\mathbf{g}\exp\left\{\int dt\sum_i[p_i(\dot\sigma_i-\mathbf{h}\xi_i+\partial V_0/\partial\sigma_i)\right.$$
$$\left.+Tp_i^2-p_i\xi_i\mathbf{f}_i-\sigma_i\xi_i\mathbf{g}_i]+\int\mathbf{f}(\mathbf{r})J_{rr'}^{-1}\mathbf{g}(\mathbf{r}')d^3r\,d^3(r')\right\}$$
$$=\int D\sigma Dp D\mathbf{f} D\mathbf{g}$$
$$\times\exp\left\{\int dt\sum_i[p_i(\dot\sigma_i-h_i+\partial V_0/\partial\sigma_i)+Tp_i^2+\sigma_il_i]\right.$$
$$\left.+\int\mathbf{f}(\mathbf{r})J_{rr'}^{-1}\mathbf{g}(\mathbf{r}')d^3r\,d^3r'\right\}. \tag{2.9}$$

In the last equation of (2.9) we have put $h_i=\mathbf{h}\xi_i+\mathbf{f}_i\xi_i$, $l_i=-\xi_i\mathbf{g}_i+\xi_i\mathbf{l}_i$. The field $\mathbf{f}(\mathbf{r}_i)\equiv\mathbf{f}_i$ has the meaning of the molecular field acting on the spin σ_i in site i. Our problem now is to integrate (2.9) with respect to the variables σ and p, i.e., to find the functional

$$\tilde Z\{l_i\}=\int D\sigma Dp\exp\left\{\int dt\sum_i[p_i(\dot\sigma_i+\partial V_0/\partial\sigma_i-h_i)\right.$$
$$\left.+Tp_i^2+\sigma_il_i]\right\}. \tag{2.10}$$

Note that $Z\{l_i\}$ is the dynamic functional for the evolution of the problem of the variable of the evolution of the variable σ_i in a potential $V_0(\sigma_i)$ and in an external field h_i. We have therefore

$$Z\{l_i\}=\exp\left\{\int dt\sum_i l_i(t)m_i(t)\right.$$
$$\left.+\frac{1}{2}\iint dt\,dt'\sum_{ij}l_i(t)D_{ij}(t')l_j(t')+\ldots\right\}, \tag{2.11}$$

where $m_i(t)=\langle\sigma_i(t)\rangle$, $D_{ij}(t-t')=\langle\langle\sigma_i(t)\sigma_j(t')\rangle\rangle$, etc. For our purposes it suffices to use only the term written out in (2.11). (The double angle brackets denote an irreducible correlator.) Recall now that we are interested in a situation in which the σ_i are Ising variables, i.e., $u\gg1$ [see (2.5)]. The dynamics of the spins σ is then described by the Glauber equations, which take in our case the very simple form

$$\dot P(\sigma_i)=\Gamma_0[-\exp(-\beta h_i\sigma_i)P(\sigma_i)+\exp(\beta h_i\sigma_i)P(-\sigma_i)],$$
$$\beta=1/T. \tag{2.12}$$

We put $P(1)=P_+$, $P(-1)=P_-$. Then $m_i\equiv m\{h_i(t)\}=P_+-P_-$ and we obtain from (2.12), recognizing that $P_++P_-=1$,

$$\dot m+2\Gamma_0m\operatorname{ch}\beta h=2\Gamma_0\operatorname{sh}\beta h. \tag{2.13}$$

We shall see presently that in the temperature range of interest to us (near the transition point) the characteristic time of variation of the field $\mathbf{f}(\mathbf{r})$, meaning also h_i, is long compared with Γ_0^{-1}. The solution of (2.13) can therefore be written in the form

$$m=\operatorname{th}\beta h-\beta\dot h/\Gamma_0. \tag{2.14}$$

We have neglected the higher derivatives of h with respect to t and the corrections to Γ_0 of second order in βh (they will be shown to be insignificant).

The quantity D_{ij} calculation with the aid of (2.13) is found to be

$$D_{ij}(t)\approx2\delta_{ij}\delta(t)/\Gamma_0,\quad|t|\gg\Gamma_0^{-1}. \tag{2.15}$$

Substituting (2.15) and (2.14) in (2.11) and (2.9) and recalling the definitions of h_i and I_i [see (2.9)] we obtain an expression for $Z\{\mathbf{l}_i\}$ in the form

$$Z\{\mathbf{l}_i\}=\int D\mathbf{f}D\mathbf{g}\exp\left\{\int dt\sum_i\left[\frac{1}{\Gamma_0}\xi_i^\alpha\xi_i^\beta g_i^\alpha g_i^\beta\right.\right.$$
$$+\mathbf{g}_i\xi_i\left(-\operatorname{th}\frac{(\mathbf{f}_i+\mathbf{h}_i)\xi_i}{T}\right.$$
$$\left.\left.+\frac{1}{T\Gamma_0}\xi_i\dot{\mathbf{f}}_i\right)\right]+\int\mathbf{g}(\mathbf{r})J_{rr'}^{-1}\mathbf{f}(\mathbf{r}')d^3r\,d^3r'$$
$$\left.+\int\mathbf{l}(\mathbf{r})J_{rr'}^{-1}\mathbf{f}(\mathbf{r}')d^3r\,d^3r'\right\}. \tag{2.16}$$

We expand the exponent in (2.16) in powers of $\mathbf{f}$ up to terms $\mathbf{f}^3$ and $\nabla^2\mathbf{f}$. (Neglect of the higher derivatives is justified by the condition $c\varkappa^{-3}\gg1$.) Equation (2.16) is reduced with the aid of (2.3) to the form

$$Z\{\mathbf{l}_i\}=\int D\mathbf{f}D\mathbf{g}\exp\left\{\int dt\,d^3r\left[\frac{1}{3\Gamma_0}\mathbf{g}^2(\mathbf{r})\right.\right.$$
$$+\frac{1}{\Gamma_0}V_{\alpha\beta}(\mathbf{r})g^\alpha(\mathbf{r})g^\beta(\mathbf{r})$$
$$+g^\alpha(\mathbf{r})\left(\tau f^\alpha(\mathbf{r})-\frac{\nabla^2}{\varkappa^2}f^\alpha(\mathbf{r})+\frac{1}{T}V_{\alpha\beta}(\mathbf{r})f^\beta(\mathbf{r})\right.$$
$$+\frac{c}{15T^3}f^\alpha\mathbf{f}^2+\frac{1}{3T\Gamma_0}\dot f^\alpha$$
$$\left.\left.\left.+\frac{1}{T\Gamma_0}V_{\alpha\beta}(\mathbf{r})\dot f^\beta(\mathbf{r})\right)\right]+\int\mathbf{l}(\mathbf{r})J_{rr'}^{-1}\mathbf{f}(\mathbf{r}')d^3r\,d^3r'\right\}. \tag{2.17}$$

Here $\tau=1-c/3T$, $V_{\alpha\beta}(\mathbf{r})=\Sigma_i\xi_i^\alpha\xi_i^\beta\delta(\mathbf{r}-\mathbf{r}_i)-c\delta^{\alpha\beta}/3$. We shall be interested hereafter in large-scale fluctuations of the field $\mathbf{f}(\mathbf{r})$, and have therefore neglected the fluctuations of the coefficient of $f^\alpha(\mathbf{f})^2$ in (2.17). $V_{\alpha\beta}(\mathbf{r})$ is a random Gaussian field having the following properties:

$$\overline{V_{\alpha\beta}(\mathbf{r})}=0,$$
$$\overline{V_{\alpha\beta}(\mathbf{r})V_{\gamma\delta}(\mathbf{r}')}=\frac{c}{15}\delta(\mathbf{r}-\mathbf{r}')(\delta_{\alpha\beta}\delta_{\gamma\delta}+\delta_{\alpha\gamma}\delta_{\beta\delta}+\delta_{\alpha\delta}\delta_{\beta\gamma}). \tag{2.18}$$

We consider hereafter the system for $|\tau|\ll 1$ ($T\approx(c/3)J_0$).

3. REGION OF FERROMAGNETIC CORRELATIONS

Consider the exponential $S(\mathbf{f},\mathbf{g})$ in (2.17) in the quadratic approximation:

$$S_0(\mathbf{f},\mathbf{g})=\int dt\,d^3r\left\{\frac{1}{3\Gamma_0}\mathbf{g}^2+g^\alpha\left[\tau f^\alpha(\mathbf{r})-\frac{\nabla^2}{\varkappa^2}f^\alpha+\frac{1}{\Gamma_0}\dot{f}^\alpha\right]\right\}. \tag{3.1}$$

The susceptibility $G_0(\mathbf{k},\omega)$ and the correlation function $K_0(\mathbf{k},\omega)$ of the field $\mathbf{f}(\mathbf{r})$, which pertain to Eq. (3.1), are of the form

$$G_0^{\alpha\beta}(\mathbf{k},\omega)=\langle g^\alpha(\mathbf{k},\omega)f^\beta(-\mathbf{k},-\omega)\rangle_0=\frac{\delta^{\alpha\beta}}{-i\omega/\Gamma_0+\tau+k^2/\varkappa^2},$$

$$K_0^{\alpha\beta}(\mathbf{k},\omega)=\langle f^\alpha(\mathbf{k},\omega)f^\beta(-\mathbf{k},-\omega)\rangle_0$$
$$=\frac{2c}{3\Gamma_0}\frac{\delta^{\alpha\beta}}{(\omega/\Gamma)^2+(\tau+k^2/\varkappa^2)^2}. \tag{3.2}$$

Expression (3.2) for the susceptibility G_0 corresponds to purely dissipative dynamics of the field $\mathbf{f}$. It is seen from (2.17) that the bare functions (3.2) require two types of correction: for the nonlinear vertex and for the random field $V_{\alpha\beta}(\mathbf{r})$. The diagrams that describe the lowest-order corrections to the susceptibility are shown in Fig. 1. The line with the arrow denotes G_0 (the arrow is directed from $\mathbf{f}$ to $\mathbf{g}$), and the line without the arrow denotes K_0; the dashed line denotes the correlator (2.18). The susceptibility calculated with allowance for these diagrams is

$$G^{-1}(\mathbf{k},\omega)=\left(\tau-0{,}1\frac{\varkappa^3}{c}\tau^{1/2}\right)+\frac{k^2}{\varkappa^2}-\frac{i\omega}{\Gamma_0}\left[1-\frac{\varkappa^3}{2c\tau^{1/2}}\frac{1}{1+(1+i\omega/\Gamma_0\tau)^{1/2}}\right]. \tag{3.3}$$

It can be seen from (3.3) that at temperatures

$$\tau\gg\tau_c\approx(\varkappa^3/2c)^2 \tag{3.4}$$

the susceptibility and the correlation functions are given in the leading approximation by Eqs. (3.2). Equation (3.3) describes in this region the corrections to $G(\mathbf{k},\omega)$ in first order in the parameter $\varkappa^3/c\tau^{1/2}$. An expression of equal accuracy for $K(\mathbf{k},\omega)$ is obtained from (3.3) with the aid of the fluctuation-dissipation theorem:

$$K(\mathbf{k},\omega)=\frac{2T}{\omega}\operatorname{Im}G(\mathbf{k},\omega). \tag{3.5}$$

It is clear from the foregoing reasoning that Eqs. (3.2)–(3.5) are valid in the frequency region $\omega/\Gamma_0\ll 1$. At $\omega/\Gamma_0\gg\tau$ the equations become somewhat simpler:

$$G^{-1}(\mathbf{k},\omega)=\tau_R+\frac{k^2}{\varkappa^2}-\frac{i\omega}{\Gamma_0}\left[1-\frac{\varkappa^3}{2c(\omega/\Gamma_0)^{1/2}}e^{-i\pi/4}\right],$$
$$K(\mathbf{k},\omega)=\frac{2c}{3\Gamma}\frac{1}{\omega^2/\Gamma_\omega^2+(\tau_R+k^2/\varkappa^2+(\omega/2\Gamma_0)^{1/2}\varkappa^3)^2},$$
$$\tau\ll\frac{\omega}{\Gamma_0}\ll 1. \tag{3.6}$$

We have put here

$$\tau_R=\tau-0{,}1\frac{\varkappa^3}{c}\tau^{1/2},\qquad \Gamma_\omega=\frac{\Gamma_0}{1-\varkappa^3/c(\omega/\Gamma_0)^{1/2}}.$$

It follows from Eqs. (3.6) that at relatively high frequencies $\omega/\Gamma_0\gtrsim(\tau/\varkappa^3)^2$ the susceptibility and the correlation function become independent of τ.

We examine now the variation of the susceptibility in the presence of an external magnetic field $\mathbf{h}$. To this end we introduce a term $-g^1h^1$ into the exponent of Eq. (2.17) for the dynamic functional and represent the field $\mathbf{f}$ in the form

$$f^\alpha(\mathbf{r})=m\delta^{1,\alpha}+\tilde{f}^\alpha(\mathbf{r}),$$

where $\langle\tilde{f}^\alpha(\mathbf{r})\rangle=0$ and m is the moment induced by the external field h. Since the random term $V_{\beta\alpha}(\mathbf{r})f^\beta(\mathbf{r})$ is small, the equation for m takes the usual form

$$\tau m+gm^3=h, \tag{3.7}$$

where $g=c/15T^3\approx 9/5c^2$. The exponent $S(\tilde{f},\mathbf{g})$ is then

$$S(\tilde{\mathbf{f}},\mathbf{g})=\int dt\,d^3r\left\{\frac{1}{3\Gamma_0}\mathbf{g}^2(\mathbf{r})+g^\alpha(\mathbf{r})\left[\tau\tilde{f}^\alpha-\frac{\nabla^2}{\varkappa^2}\tilde{f}^\alpha-\frac{3}{c}V_{\alpha\beta}\tilde{f}^\beta+g\tilde{f}^\alpha\tilde{\mathbf{f}}^2-\frac{3}{c}mV^{\alpha 1}+g(2m^2\tilde{f}^1\delta^{1,\alpha}+m\tilde{f}^2\delta^{1,\alpha}+m^2\tilde{f}^\alpha+2m\tilde{f}^\alpha\tilde{f}^1)\right]\right\}. \tag{3.8}$$

(We have written out only the terms of importance for the exposition that follows.) As seen from (3.8), the external magnetic field h has led to the appearance of a new triple vertex of type $\mathbf{g}\tilde{\mathbf{f}}^2$ and of a random field $\tilde{\mathbf{h}}^\alpha\propto mV^{\alpha 1}$, which acts on $\tilde{\mathbf{f}}(\mathbf{r})$. The random field is known to add to the static correlation function a contribution proportional to the square of the Lorentzian (see Ref. 8). The correction to the susceptibility, necessitated by the new vertex, is illustrated by the diagram in Fig. 2. We introduce the longitudinal and transverse susceptibilities in accordance with the equation

$$G^{\alpha\beta}(\mathbf{k},\omega)=G_\parallel(\mathbf{k},\omega)\delta^{1\alpha}\delta^{1\beta}+G_\perp(\mathbf{k},\omega)(\delta^{\alpha\beta}-\delta^{1\alpha}\delta^{1\beta}). \tag{3.9}$$

Calculation of the diagram of Fig. 2 leads to the following expressions for $\sigma G_\parallel^{-1}$ and $\sigma G_\perp^{-1}$ at $k=0$:

$$\delta G_\parallel^{-1}(0,\omega)\approx 10\frac{m^2}{c^2}\frac{\varkappa^3}{c\tau_\parallel^{1/2}}\frac{1}{1+[1-i\omega/(\Gamma_0\tau_\parallel)]^{1/2}},$$

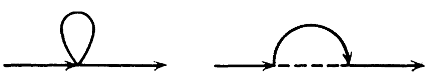

FIG. 1.

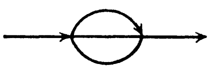

FIG. 2.

$$\delta G_\perp^{-1}(0,\omega) \approx \frac{m^2}{c^2}\frac{\varkappa^3}{c\tau_\perp^{1/2}}\frac{1}{1+[1-i\omega/(\Gamma_0\tau_\perp)]^{1/2}}, \quad (3.10)$$

where $\tau_\parallel = \tau + 27m^2/5c^2$, $\tau_\perp = \tau + 9m^2/5c^2$. Separating the term $\sigma G^{-1}(0,0)$ that renormalizes τ, and taking (3.3) into account, we get

$$G_\parallel^{-1}(0,\omega) = \tau_\parallel - \frac{i\omega}{\Gamma_0}\left\{1 - \frac{\varkappa^3}{2c\tau_\parallel^{1/2}}\frac{1}{1+(1+i\omega/\Gamma_0\tau_\parallel)^{1/2}} - \frac{5m^2}{\tau_\parallel c^2}\frac{\varkappa^3}{c\tau_\parallel^{1/2}}\frac{1}{[1+(1+i\omega/\Gamma_0\tau_\parallel)^{1/2}]^2}\right\}, \quad (3.11)$$

$$G_\perp^{-1}(0,\omega) = \tau_\perp - \frac{i\omega}{\Gamma_0}\left\{1 - \frac{\varkappa^3}{2c\tau_\perp^{1/2}}\frac{1}{1+(1+i\omega/\Gamma_0\tau_\perp)^{1/2}} - \frac{m^2}{2\tau_\perp c^2}\frac{\varkappa^3}{c\tau_\perp^{1/2}}\frac{1}{[1+(1+i\omega/\Gamma_0\tau_\perp)^{1/2}]^2}\right\}.$$

The last term in (3.11) is always small compared with unity. We see thus that application of an external magnetic field gives rise only to small contributions to the susceptibility and to the correlation function. The paramagnetic resonance typical of ferromagnet is absent in this case. The reason is that in view of the strong anisotropy there is no spin precession in the magnetic field.

4. CONCLUSION

We have considered the dynamics of amorphous magnets with strong random anisotropy in the ferromagnetic-correlation region $\tau = (T - T_1)/T_1 \gg \tau_c$, where $T_1 = 1/3c\int J(r)d^3r$ (c is the density of the magnetic ions), $\tau_c \approx (\varkappa^3/2c)^2$, $\varkappa^{-1}$ is the radius along which the exchange $J(r)$ interaction falls off. It was assumed here that this radius is large compared with the interatomic distance: $c\varkappa^{-3} \gg 1$. [At $c \approx 0.5a^{-3}$, where a is the interatomic distance, it actually suffices to have $\varkappa^{-1} \approx (1.5$ to $2)a$, which is quite realistic.] As shown in Ref. 8, in this temperature region the static correlation function $G(\mathbf{p})$ takes the form $G(\mathbf{p}) \propto 1/(\tau + p^2/\varkappa^2)$ typical of a pure Heisenberg ferromagnet. In the present paper we have calculated the susceptibility $G(\mathbf{k},\omega)$ and the correlation function $K(\mathbf{k},\omega)$. It was found that in the leading approximation the dynamics is purely dissipative,

$$G(\mathbf{k},\omega) \propto 1/[(i\omega/\Gamma_0) + \tau + k^2/\varkappa^2],$$

$$K(\mathbf{k},\omega) \propto 1/[(\omega/\Gamma_0)^2 + (\tau + k^2/\varkappa^2)^2],$$

where Γ_0 is the characteristic spin-flip frequency.

Thus, the analogy with Heisenberg ferromagnets is lost in the dynamics. The reason is that owing to the strong anisotropy the total spin is not conserved in the system. Note that the dissipative dynamics was derived directly from a microscopic model, and that the dynamic-functional method[9] was used. The first-order corrections to $G(\mathbf{k},\omega)$ and $K(\mathbf{k},\omega)$ with respect to the parameter τ/τ_c were also calculated [see Eqs. (3.3)–(3.6)].

An external magnetic field gives rise to additional corrections to the susceptibility $G(\mathbf{k},\omega)$, which we have calculated for $\mathbf{k} = 0$ [see (3.1) and (3.11)]. Since the strong anisotropy prevents spin precession in the magnetic field, there is no paramagnetic resonance in this case

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