Fast coherent relaxation in a ferromagnet nanoparticle assembly

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(Received 2 March 2011; revised manuscript received 10 July 2011; published 5 October 2011)

The idea to establish coherent relaxation in an assembly of magnetic moments by placing it inside a passive resonator is applied to the case of single-domain ferromagnetic nanoparticles. The dynamics of magnetization inversion of a nanoparticle is governed by the Landau–Lifshitz (spin-lattice) relaxation and radiation damping. The numeric simulations enable us to account for the interparticle dipole-dipole interactions in a rigorous way. A new mechanism of fast relaxation in a nanoparticle assembly is suggested and described.

DOI: 10.1103/PhysRevB.84.144412 PACS number(s): 75.75.Jn, 76.60.−k, 76.20.+q, 76.50.+g

I. INTRODUCTION

Fast magnetic switching is one of the main goals in magnetic recording. In a medium consisting of fine grains, like FePt $L1_0$, the problem is to accelerate as much as possible the switching process for a single particle with uniaxial anisotropy. Besides simple application of an inverse field (i.e., the one antiparallel to the actual equilibrium direction of the magnetic moment), several other strategies were proposed, such as precessional[1–5] or heat-assisted[6,7] switching.

In the act of switching, the field pulse imparts to the magnetic moment a portion of energy sufficient to surmount the anisotropy barrier. However, as soon as the barrier is traversed, the same energy should be evacuated. If the dissipation is too slow and the magnetic moment does not get rid of the energy excess at the timescale of the field pulse, the final orientation state of the magnetic moment becomes uncertain due to a substantial probability of the return to the initial state. In other words, to ensure fast switching, a minimal damping time is necessary.

In this paper, we study the situation in which the after-switching energy dump of the magnetic moment of a single-domain particle (or the particle assembly) is facilitated by coupling of the sample to a passive electric circuit. As far as we know, the problem of such a type was first considered by Bloembergen and Pound[8] with respect to the magnetic resonance technique, since the latter essentially involves inductive coupling of the magnetic moment to the pick-up coil. This type of relaxation was then termed *radiation damping*. The latter effect manifests the stronger, the greater the number of particles in the assembly. The cause is that the feedback field, induced in the resonator by the joint action of all the members of the assembly, synchronizes their individual motions and accelerates the inversion process. Under these conditions, the energy discharge (radiation) pulse shortens, while its power enhances proportionally. A similar phenomenon is well known in optics as *superradiation* and is well studied. The universal signature of this effect is that the damping rate becomes approximately proportional to the number of spins participating in the collective relaxation.[9]

Here, we develop the radiation damping (RD) idea for the case of ferromagnetic particles, noting, however, that along with apparently similar features, the physics of the collective relaxation in a spin system (magnetic field, UHF wave range) differs significantly from its optical analogue (electric field, optical wave range). The RD problem for nuclear and electronic spins in paramagnets was investigated theoretically in detail in a number of recent papers,[10–17] and the experimental verifications for nuclear spins on both solid and liquid samples have been obtained.[18–20] In the present paper, we apply the RD concept to an assembly of single-domain ferromagnetic particles, thus following the line set out in paper [21] in which the Landau–Lifshitz equation for macroscopic magnetization of a sample was used to study the collective magnetic relaxation in nanomagnets. The specificity of the problem is twofold. First, the "elementary" magnetic moments are huge with respect to the atomic scale (10^4–10^5 Bohr magnetons), thus yielding a strong effect that would not be smeared out at ambient temperatures. Second, unlike paramagnets, here the switching process cannot be split into a fast transverse and slow longitudinal stages because in a ferromagnet, the spin-spin and spin-lattice relaxation rates practically coincide. Formally, this means replacing the Bloch equation by the Landau–Lifshitz equation. Therefore, the RD effect in a ferromagnet nanoparticle assembly is a unique phenomenon. As our analysis shows, it is able to establish a new mechanism of substantial reduction of the magnetic switching time.

Consider an assembly of single-domain particles with uniaxial anisotropy, coupled to each other by the magnetic dipole-dipole interaction. For simplicity, the assembly is assumed to be "orientationally textured"; that is, the particles are embedded in a solid matrix in such a way that all their easy axes are parallel. The initial state can be prepared by subjecting the system to a strong (compared with the particle coercive force) field pulse that aligns all the magnetic moments in the same direction. After the field is turned off, they rest in the respective minima of the anisotropy potential (with allowance for the effect of the dipolar field). Therefore, the system is magnetized almost to saturation and because the anisotropy energy is assumed to be much higher than thermal energy, this magnetic configuration is stable. When the system is subjected to a stepwise field directed opposite to the initial orientation, the magnetization begins to relax toward the new equilibrium. When turning, the magnetic moments change the magnetic flux through the coil, which is a part of a passive *LCR* circuit. The electric current induced in the coil generates a magnetic

field that acts on all the moving magnetic moments of the assembly. This feedback effect redistributes the energy and its losses in the system and thus modifies the magnetic relaxation.

II. HIGH-FREQUENCY MAGNETODYNAMICS OF SINGLE-DOMAIN PARTICLES IN A RESONATOR. BASIC SET OF EQUATIONS

We describe the magnetodynamics of the particle assembly with the aid of the standard Landau–Lifshitz equation, so that for the kth particle (i.e., for the magnetic moment $\mu^{(k)}$), one has

$$\frac{d\boldsymbol{\mu}^{(k)}}{dt} = -|\gamma|(\boldsymbol{\mu}^{(k)} \times \boldsymbol{\mathcal{H}}^{(k)}) - \frac{\alpha|\gamma|}{\mu}(\boldsymbol{\mu}^{(k)} \times (\boldsymbol{\mu}^{(k)} \times \boldsymbol{\mathcal{H}}^{(k)})); \tag{1}$$

here, α is the dimensionless relaxation parameter, and γ is the gyromagnetic ratio for electrons. The field $\boldsymbol{\mathcal{H}}^{(k)}$ in Eq. (1) is a full magnetic field acting on the kth particle. In the considered case, it comprises:

(1) the external constant field $\boldsymbol{H}_0$ along the Oz axis;
(2) the uniaxial anisotropy field

$$\boldsymbol{H}_A = (H_A/\mu)(\boldsymbol{\mu} \cdot \boldsymbol{n})\boldsymbol{n}, \quad H_A = 2E_A/\mu, \tag{2}$$

where $\boldsymbol{n}$ is the unit vector of the easy axis and E_A is the particle anisotropy energy;

(3) the feedback field $\boldsymbol{H} = (H,0,0)$ generated by the current induced in the coil, whose axis is directed along Ox; and
(4) the dipolar magnetic field $\boldsymbol{H}_d^{(k)}$ caused by the interparticle dipole-dipole pair interactions.

Setting Oz along the common direction of the particle easy axis, one has $\boldsymbol{n} = (0,0,1)$, so that the effective field takes the form

$$\boldsymbol{\mathcal{H}} = (H + H_{dx}, H_{dy}, H_0 + \mu_z H_A/\mu + H_{dz}). \tag{3}$$

The local dipolar magnetic field $\boldsymbol{H}_d^{(k)} = -\partial U_{dd}/\partial \boldsymbol{\mu}^{(k)}$ at the site of the kth particle is defined from the sum of the pairwise dipole-dipole energy contributions:

$$U_{dd} = \sum_{\substack{k,m \\ k>m}}^{N} \left[\frac{1}{r_{km}^3}(\boldsymbol{\mu}^{(k)}\boldsymbol{\mu}^{(m)}) - \frac{3}{\boldsymbol{r}_{km}^5}(\boldsymbol{\mu}^{(k)}\boldsymbol{r}_{km})(\boldsymbol{\mu}^{(m)}\boldsymbol{r}_{km}) \right], \tag{4}$$

where $\boldsymbol{r}_{km}$ is the radius vector connecting the kth and mth particles, and N is the number of particles.

We define the reference frequency of precession in the external field $\boldsymbol{H}_0$ and those related to the feedback, dipolar, and anisotropy fields, respectively as:

$$\omega_0 = |\gamma|H_0, \ \ \omega_H = |\gamma|H, \ \ \omega_d = |\gamma|\mu/a^3, \ \ \omega_A = |\gamma|H_A. \tag{5}$$

Substituting Eqs. (2)–(5) into Eq. (1), one obtains

$$\begin{aligned}
\dot{\mu}_x^{(k)} &= -\left(\omega_0 + \omega_A\frac{\mu_z^{(k)}}{\mu}\right)\mu_y^{(k)} - \gamma\left(\mu_y^{(k)}H_{dz}^{(k)} - \mu_z^{(k)}H_{dy}^{(k)}\right) + \alpha\gamma(H + H_{dx})\frac{\mu_y^{(k)2} + \mu_z^{(k)2}}{\mu} - \alpha\left(\omega_0 + \gamma H_{dz} + \omega_A\frac{\mu_z^{(k)}}{\mu}\right) \\
&\quad \times \frac{\mu_x^{(k)}\mu_z^{(k)}}{\mu} - \frac{\alpha\gamma H_{dy}}{\mu}\mu_x\mu_y, \\
\dot{\mu}_y^{(k)} &= \left(\omega_0 + \omega_A\frac{\mu_z^{(k)}}{\mu}\right)\mu_x^{(k)} - \gamma\mu_z^{(k)}H - \gamma\left(\mu_z^{(k)}H_{dx}^{(k)} - \mu_x^{(k)}H_{dz}^{(k)}\right) - \alpha\gamma(H + H_{dx})\frac{\mu_x^{(k)}\mu_y^{(k)}}{\mu} - \alpha\left(\omega_0 + \gamma H_{dz} + \omega_A\frac{\mu_z^{(k)}}{\mu}\right) \\
&\quad \times \frac{\mu_y^{(k)}\mu_z^{(k)}}{\mu} + \frac{\alpha\gamma H_{dy}}{\mu}\left(\mu_x^{(k)2} + \mu_z^{(k)2}\right), \\
\dot{\mu}_z^{(k)} &= \gamma\mu_y^{(k)}H - \gamma\left(\mu_x^{(k)}H_{dy}^{(k)} - \mu_y^{(k)}H_{dx}^{(k)}\right) - \alpha\gamma(H + H_{dx})\frac{\mu_x^{(k)}\mu_z^{(k)}}{\mu} + \alpha\left(\omega_0 + \gamma H_{dz} + \omega_A\frac{\mu_z^{(k)}}{\mu}\right) \\
&\quad \times \frac{\mu_x^{(k)2} + \mu_y^{(k)2}}{\mu} - \frac{\alpha\gamma H_{dy}}{\mu}\mu_y^{(k)}\mu_z^{(k)}.
\end{aligned} \tag{6}$$

We remark that in Eqs. (6), two "collective" factors are present, whose magnitudes are determined by joint contributions of all the particles of the assembly: the feedback field H, which is proportional to the rate of change of the full magnetic moment of the assembly, and the dipolar field H_d.

Introducing dimensionless time

$$\tilde{t} = \omega_0 t, \tag{7}$$

and scaling the reference frequencies of Eq. (5) with the Larmor value ω_0, one gets a set of dimensionless parameters,

$$p_H = \frac{\omega_H}{\omega_0} = \frac{H}{H_0}, \quad p_d = \frac{\omega_d}{\omega_0} = \frac{\mu}{a^3 H_0}, \quad p_A = \frac{\omega_A}{\omega_0} = \frac{H_A}{H_0}, \tag{8}$$

where a is the mean interparticle distance. Moreover, because Eq. (1) conserves the modulus of the magnetic moment, it is convenient to introduce unit vectors $\boldsymbol{e}^{(k)} = \boldsymbol{\mu}^{(k)}/\mu$ that denote the respective directions. As a result, Eqs. (6) transform to

$$\begin{aligned}
\dot{e}_x^{(k)} &= -\left(1 + p_A e_z^{(k)}\right)e_y^{(k)} - p_d\left(e_y^{(k)}\tilde{H}_{dz}^{(k)} - e_z^{(k)}\tilde{H}_{dy}^{(k)}\right) \\
&\quad + \alpha(p_H + p_d\tilde{H}_{dx})\left(e_y^{(k)2} + e_z^{(k)2}\right) \\
&\quad - \alpha\left(1 + p_A e_z^{(k)} + p_d\tilde{H}_{dz}\right)e_x^{(k)}e_z^{(k)} - \alpha p_d\tilde{H}_{dy}e_x^{(k)}e_y^{(k)}, \\
\dot{e}_y^{(k)} &= \left(1 + p_A e_z^{(k)}\right)e_x^{(k)} - p_H e_z^{(k)} - p_d\left(e_z^{(k)}\tilde{H}_{dx}^{(k)} - e_x^{(k)}\tilde{H}_{dz}^{(k)}\right) \\
&\quad - \alpha(p_H + p_d\tilde{H}_{dx})e_x^{(k)}e_y^{(k)} - \alpha\left(1 + p_d\tilde{H}_{dz} + p_A e_z^{(k)}\right) \\
&\quad \times e_y^{(k)}e_z^{(k)} + \alpha p_d\tilde{H}_{dy}\left(e_x^{(k)2} + e_z^{(k)2}\right),
\end{aligned}$$

$$\dot{e}_z^{(k)} = p_H e_y^{(k)} - p_d\left(e_x^{(k)}\tilde{H}_{dy}^{(k)} - e_y^{(k)}\tilde{H}_{dx}^{(k)}\right) - \alpha(p_H + p_d\tilde{H}_{dx}) \times e_x^{(k)}e_z^{(k)} + \alpha\left(1 + p_A e_z^{(k)} + p_d\tilde{H}_{dz}\right)\left(e_x^{(k)2} + e_y^{(k)2}\right) - \alpha p_d\tilde{H}_{dy}e_y^{(k)}e_z^{(k)}. \quad (9)$$

Here, time derivatives are defined with respect to time $\tilde{t}$, and the dimensionless dipolar field in the kth site is

$$\boldsymbol{H}_d^{(k)}/H_0 = p_d\tilde{\boldsymbol{H}}_d^{(k)}, \qquad \tilde{\boldsymbol{H}}_d^{(k)} = \sum_{\substack{m=1\\ m\neq k}}^{N}\left[\frac{3}{\tilde{r}_{km}^5}\tilde{\boldsymbol{r}}_{km}(\boldsymbol{e}^{(m)}\tilde{\boldsymbol{r}}_{km}) - \frac{1}{\tilde{r}_{km}^3}\boldsymbol{e}^{(m)}\right], \quad (10)$$

where $\tilde{\boldsymbol{r}}_{lm} = \boldsymbol{r}_{lm}/a$ are dimensionless vectors of the interparticle distances.

First, we consider Eqs. (9) for a single particle and assume that the feedback and dipolar fields are negligible ($H = H_d = 0$). Thence, these equations uncouple and form independent triads, each subset referring to a separate particle:

$$\dot{e}_x = -(\omega_0 + \omega_A e_z)[e_y + \alpha e_x e_z], \quad \dot{e}_y = (\omega_0 + \omega_A e_z)[e_x - \alpha e_y e_z], \quad \dot{e}_z = \alpha(\omega_0 + \omega_A e_z)\left[1 - e_z^2\right]. \quad (11)$$

Moreover, in each triad, the equation for e_z, the projection of the magnetic moment on the field $\boldsymbol{H}_0$, is closed, thus yielding a single equation to describe the magnetization inversion (switching). Although this equation admits a solution in quadratures, it turns out to be easier to perform its numerical integration. For that, we take the initial condition in the form $e_z(t=0) = -1 + \delta$, where δ is a positive parameter necessary to deviate the magnetic moment from the unstable equilibrium state $\boldsymbol{e} \parallel -\boldsymbol{H}_0$. The solution obtained shows that the magnetic moment inversion develops monotonically with reference time $\tau_{\rm inv} = 1/\alpha(\omega_0 + \omega_A)$, so that for $t \gg \tau_{\rm inv}$ the projection $e_z(t)$ turns to unity. Such a regime takes place, however, only for $\omega_0 > \omega_A$—that is, when the remagnetizing field exceeds the particle coercive force. In the opposite case, $\omega_0 < \omega_A$, the magnetic moment does not possess enough energy to surmount the potential barrier, and, instead of inversion, relaxation would occur to the same minimum: $e_z(t \gg \tau_{\rm inv}) \to -1$. We remark that the considerations given above are valid only for the athermic case (i.e., zero temperature). At finite temperatures, the magnetic inversion due to superparamagnetism can occur even at "subthreshold" fields of the order $(\omega_A - \omega_0)/\gamma \sim kT/\mu$.

III. FEEDBACK FIELD EQUATION

As mentioned, the Ox axis of the coordinate frame is directed along the axis of the induction coil of the LCR circuit. Therefore, the nonstationary electromotive force and, accordingly, the electric current I in the circuit are due to the time change of the x-component of the net magnetic moment of the system. The corresponding Kirchhoff equation is

$$L\frac{dI}{dt} + RI + \frac{1}{C}\int_0^t I(t')dt' = -\frac{d\Phi}{dt}, \quad \Phi = (4\pi/c)n\eta A m_x, \quad (12)$$

where Φ is the magnetic flux in a coil with n turns and the cross-section area A; the constant c is the speed of light in vacuum, $\eta = V/V_c$ is the coil filling factor, V is the volume of the sample containing ferromagnetic particles, and V_c is the inner volume of the coil. The quantity

$$m_x = (\mu/V)\sum_l e_x^{(l)} \quad (13)$$

is the x component of the sample magnetization.

The induced current generates in the coil the magnetic (feedback) field

$$H = (4\pi n/cl)I. \quad (14)$$

The self-induction coefficient of a coil of length l in the simplest approximation is $L = 4\pi n^2 A/lc^2$. Differentiating Eq. (12) with respect to time, using the time scale $\tilde{t}$ and the variable p_H, one arrives at the equation

$$\frac{d^2}{d\tilde{t}^2}p_H + 2\frac{\gamma_r}{\omega_0}\frac{d}{d\tilde{t}}p_H + \left(\frac{\omega_r}{\omega_0}\right)^2 p_H = -4\pi\beta\left(\frac{1}{N}\frac{d^2}{d\tilde{t}^2}\sum_{l=1}^{N}e_x^{(l)}\right), \quad (15)$$

where the coefficients in the left-hand side are expressed in terms of the LCR circuit parameters as $2\gamma_r = R/L = \omega_r/Q, \omega_r = 1/\sqrt{LC}$, with Q being the quality factor. The function in parentheses in the right-hand side of Eq. (15) is the averaged second derivative of the x component of the unit vector of magnetization, whereas the coefficient

$$\beta = \eta N\mu/(VH_0) \quad (16)$$

alongside it, determines the intensity of inductive coupling between the particle assembly and the coil. Using estimation $a \approx (V/N)^{1/3}$ and the second of the definitions in Eqs. (8), parameter β may be presented in the form ηp_d.

We note that the feedback field H, induced by the resonator, is the governing factor of the inversion process. It emerges as soon as the system of interacting magnetic moments, having been put in a nonequilibrium state at $t = 0$, begins to move toward equilibrium. Because the feedback field unites the contributions from all particles, one should expect that this factor would reduce noncoherent contributions to the dynamic magnetization $N\mu\langle e\rangle$, thus making the inversion process more coherent.

IV. MAGNETIZATION INVERSION UNDER FEEDBACK

The set of $3N$ Eqs. (9) together with Eq. (15) is solved numerically for a sample containing N particles. We performed simulations for different numbers of spins and have checked that the results are very similar for N from 18 to 125.

For a given value of initial polarization $e_z(0)$, a variety of initial orientations of individual vectors $e_z^{(k)}(0)$ are possible. These orientations are obtained with a Monte Carlo–like technique. A random configuration of vectors $\{e_z^{(k)\rm init}(0)\}$ is taken, and the corresponding total polarization $e_z^{\rm init}(0) = N^{-1}\sum_k e_z^{(k)\rm init}(0)$ is evaluated. A new direction is chosen randomly for each spin, and the new total polarization is calculated. If it is less than the initial value, the array with the changed magnetic moment direction distribution is chosen as the second iteration; otherwise, it is rejected. This procedure

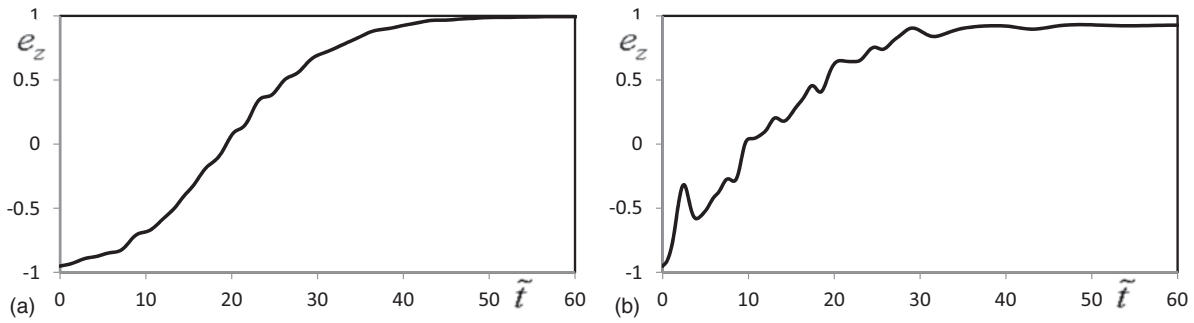

FIG. 1. Magnetization evolution under inverted field in the presence of dipole interactions for the precession damping parameter $\alpha = 0.1$, initial polarization $e_z(0) = -0.95$, and $p_A = 0$; (a) $p_d = 0.07$, (b) $p_d = 0.25$.

is repeated until the system achieves polarization $e_z(0)$, and the obtained state is taken for the initial magnetic moment distribution of the system. The details of this preliminary procedure are described in [11]. The initial conditions for Eq. (15) are zero: $p_H(0) = 0$, $\dot{p}_H(0) = 0$.

To get a notion of the particular situation in question, it is instructive to have some numerical estimations. Let the particles be made of a ferrite with the saturation magnetization $M_S = 400$ G and have the mean size $d \simeq 10$ nm. Then, the particle magnetic moment is $\mu \simeq 2 \times 10^{-16}$ emu, which is about 10^4 Bohr magnetons. Assuming that the volume content of the particles in the sample is $\phi = 10\%$, for the mean interparticle distance we obtain $a \sim d/\phi^{1/3} \sim 2d$. For the magnetizing field of a typical strength $H_0 \simeq 3300$ Oe, the Larmor frequency is $\omega_0/2\pi = \gamma H_0/2\pi \approx 10$ GHz, whereas for the dipolar field parameter, one gets $p_d = \mu/a^3 H_0 \sim \phi M_S/H_0 \sim 10^{-2}$. Setting the filling factor of the coil $\eta = 1$, we see that definitions of p_d from Eq. (8) and that of β from Eq. (16) coincide, so for the problem under study, $\beta \simeq p_d$. For the precession damping parameter, we assume a typical value $\alpha \sim 0.1$.

V. RESULTS AND DISCUSSION

The figures below demonstrate the time behavior of polarization $e_z(t) = (1/N)\sum_l e_z^{(l)}(t)$ and the feedback field H (in units of H_0). All the figures are simulated for the number of spins $N = 100$ making a sample with the geometry of an oblate bar with the dimensions $N_x = N_y = 5, N_z = 4$.

We first consider the case in which the resonator is absent ($\beta = 0$) and so is the feedback field. Figure 1 shows the regime of the pure Landau–Lifshitz relaxation in an assembly of isotropic particles coupled by the magnetic dipolar interaction. One can see that the interparticle interaction accelerates relaxation, especially at its initial stage. We note that as the value of the damping parameter α grows, the relaxation rate increases.

Typical results, which combine Landau–Lifshitz and radiation damping, are shown in Figs. 2–5. The resonator eigenfrequency is chosen to be tuned to the Larmor frequency (i.e., $\omega_r = \omega_0$; $\omega_r = \langle\omega_0^{(k)}\rangle$ in Fig. 5), and the quality factor in all the calculations is set to $Q = 10$. As seen, the relaxation process under these conditions is drastically different from that in a pure Landau–Lifshitz regime. The most important distinction is that, in the presence of the feedback (resonator), the process becomes coherent and, therefore, faster. Indeed, the magnetization practically accomplishes a complete flip during just a few precession periods ω_0^{-1}. The feedback field worked out by the resonator attains its maximal value when total magnetization passes through zero. As seen, this field can reach quite high values (e.g., exceeds the external constant field H_0, as in Fig. 2a). Also noteworthy is that, if radiation relaxation is considered by itself (no Landau–Lifshitz process), the relaxation of e_z becomes oscillatory with a substantial amplitude, and the polarization oscillations are multiple. The Landau–Lifshitz relaxation term in the magnetodynamic equations weakens these oscillations, whose traces are still visible, however (Fig. 2), and clearly fade out with the growth of α.

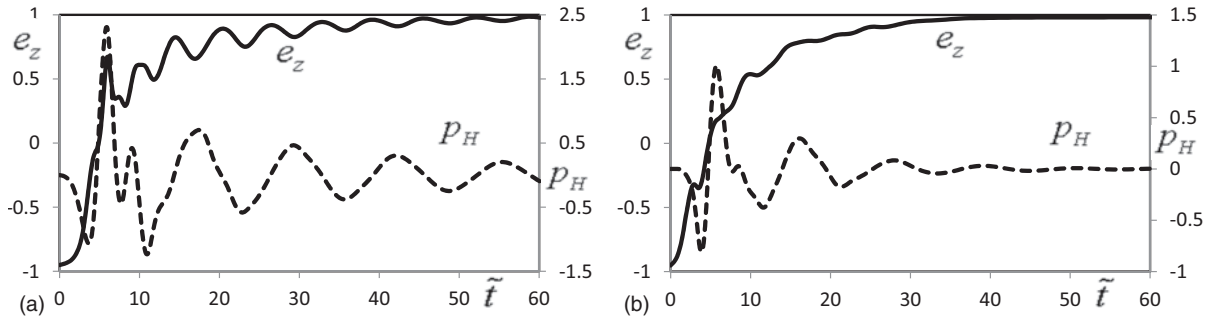

FIG. 2. Polarization e_z (left scale) and feedback field p_H (right scale) versus $\tilde{t}$ for $e_z(0) = -0.95, p_A = 0$, $Q = 10$, $\omega_r = \omega_0$. The dipole parameter $p_d = 0.2$, and parameter β has the same value; (a) $\alpha = 0.1$, (b) $\alpha = 0.2$.

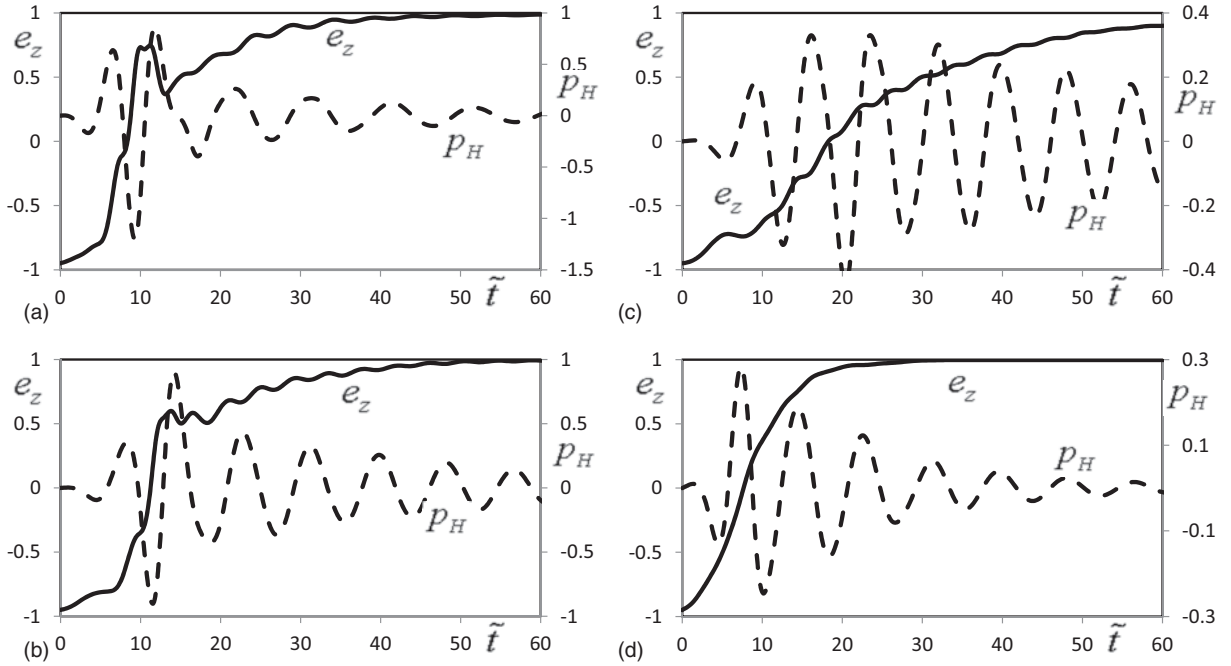

FIG. 3. Polarization e_z (left scale) and feedback field p_H (right scale) versus $\tilde{t}$ for $e_z(0) = -0.95$, $Q = 10$, $\omega_r = \omega_0$. The dipole parameter $p_d = 0.07$, and parameter β has the same value. (a) $p_A = 0$, $\alpha = 0.1$; (b) $p_A = 0.5$, $\alpha = 0.1$; (c) $p_A = 0.8$, $\alpha = 0.1$; (d) $p_A = 0.5$, $\alpha = 0.3$.

Figure 3 demonstrates the role of anisotropy and the Landau–Lifshitz damping parameter; it is seen that an increase in anisotropy makes the relaxation slower, whereas augmentation of α accelerates it.

As follows from Eqs. (9) and (15), the parameters α and β contribute to the relaxation mechanisms in a nonlinear way. To analyze these dependencies, we define the reference time τ as the interval, at the end of which polarization e_z changes its sign from negative to positive. The results are shown in Fig. 4. For small values of β and large values of α, the Landau–Lifshitz relaxation dominates. The increase of coupling between the magnetic moments and the coil leads to practically proportional growth of the radiation damping rate. Because the coupling parameter is directly proportional to N, the RD rate is proportional to the number of magnetic moments involved in the coherent relaxation. For large values of β, the RD mechanism dominates, and the relaxation rate loses its dependence on α.

For the considered RD modulation of magnetic relaxation to be observable, it should exceed the ever-present smearing of the resonance frequency ω_0. Two sources are most probable for this smearing. The first stems from the distribution of the internal demagnetizing fields caused by some possible nonsphericity of the particles. The second is due to the contribution of surface effects to the internal magnetic anisotropy fields.

To estimate the allowable extent of particle nonsphericity, we assume that, instead of a sphere, the particle shape is an ellipsoid of revolution with aspect ratio b/a. The magnetostatic term in the effective anisotropy then takes the form

$$K_{ns} = 2\pi M_s^2 \Delta N,$$

where ΔN is the difference of demagnetizing factors along the principal axes. When the particle anisometricity is not too large, the corresponding contribution to the particle internal field is

$$H_{ns} = 2K_{ns}/M_s = \frac{4\pi}{5} M_s \varepsilon^2,$$

with ε being the ellipsoid eccentricity (e.g., see, Ref. 22). Setting the aspect ratio $b/a \simeq 0.9$, for $M_s = 400$ G, we get $\varepsilon^2 \simeq 0.2$ and, hence, $H_{ns} \simeq 200$ Oe. Comparing the latter to $H_0 \simeq 3300$ Oe, one finds that the resulting spread of Larmor frequencies, $p_{ns} = \Delta\omega_0/\omega_0 = H_{ns}/H_0$, amounts to about 6%. As Figs. 1 and 2 show, the effect of the dipole interactions becomes important when $p_d \simeq 0.2$ or higher. Comparing p_{ns} and p_d, one concludes that for the considered systems, the

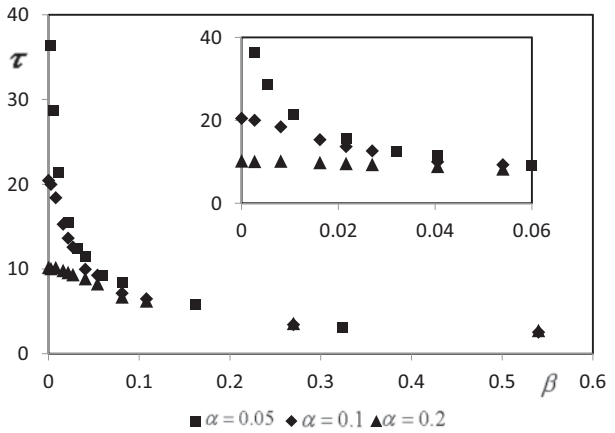

FIG. 4. Time τ versus coupling constant β; the parameters are $e_z(0) = -0.95$, $Q = 10$, $\omega_r = \omega_0$, $p_A = 0$; the values of α are indicated within the figure. The insert shows the same figure for small values of β.

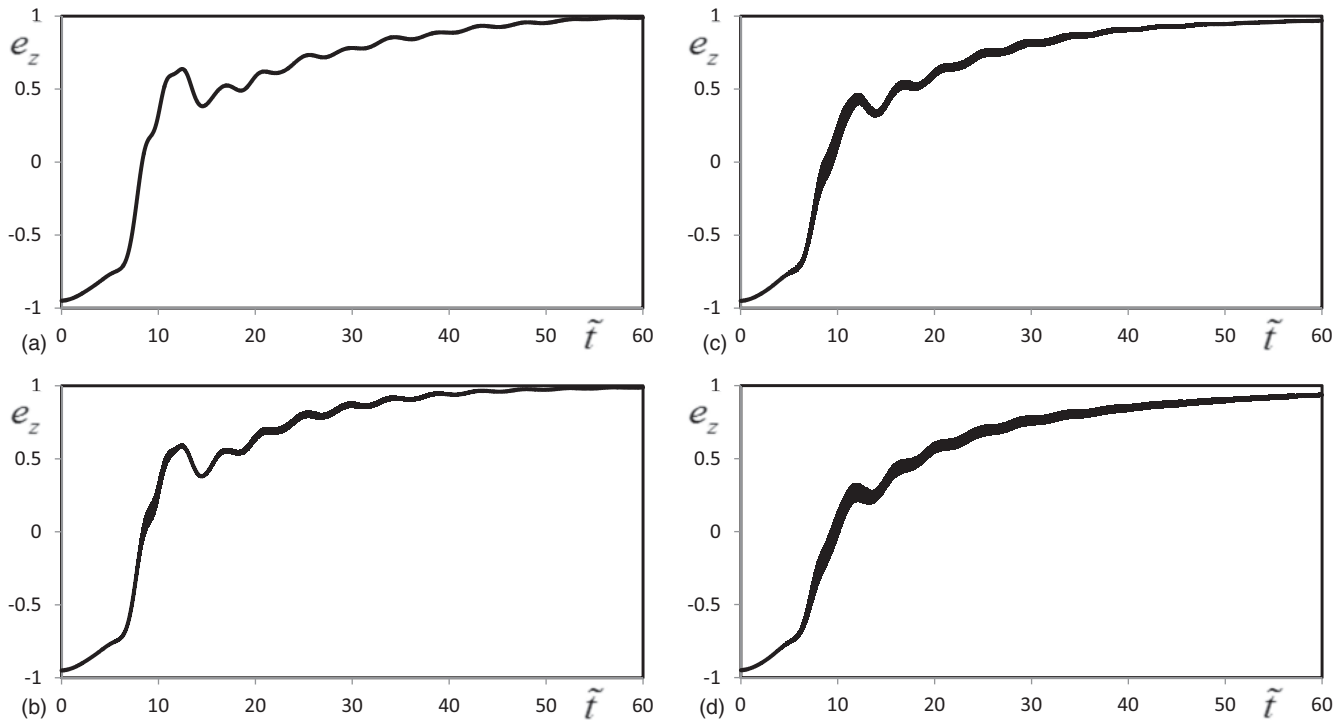

FIG. 5. Polarization e_z versus time $\tilde{t}$ for different dispersion values of the Zeeman frequency. Values of the parameters are $e_z(0) = -0.95$, $Q = 10$,$p_d = \beta = 0.07$, $p_A = 0.3$, $\alpha = 0.1$. (a) Zero dispersion; (b) averaging with dispersion $\Delta = 10\%$; (c) averaging with dispersion $\Delta = 20\%$; (d) averaging with dispersion $\Delta = 30\%$.

shape polydispersity of the particles up to 10% does not affect the results substantially.

The surface-induced spread of magnetocrystalline anisotropy of individual particles can induce much larger spread of Larmor frequencies. The results of model treatment of such a situation are shown in Fig. 5, where the RD effect is presented for a sample, where Larmor frequencies $\omega_0^{(k)}$ ($k = 1, \ldots N$) are distributed normally with the dispersions $\Delta = 10$, 20, and 30% expressed in units of the mean value of $\omega_0^{(k)}$. Each curve in Fig. 5(b)–5(d) is obtained by averaging over 60 such distributions [with the same set of initial orientations of individual vectors $e_z^{(k)}(0)$]. The width of the curves (their thickness) is obtained assuming the confidence coefficient 0.9 calculated according to Student's criterion. As seen, the increase of the dispersion, quite expectedly, entails some growth of the relaxation time, but the RD effect caused by a strong interaction of the magnetic particles with the coil sustains quite well nevertheless.

VI. CONCLUSIONS

A new mechanism of fast magnetic relaxation in a ferromagnet nanoparticle system is proposed. Being based on the theory of superradiation, it establishes that the reference time of coherent relaxation, under some conditions, is inversely proportional to the number of magnetic moments. A conceivable way to obtain such an effect is to place an assembly of magnetic moments in a resonance-tuned coil. The cooperative feedback magnetic field from the coil acting on the magnetic moments establishes the regime of coherent relaxation from the initial disequilibrium state. Such a relaxation process is very different from that caused by a pure Landau–Lifshitz relaxation. The most important issue is that the coherent process is substantially faster, so that the magnetization practically accomplishes a complete flip during just few Larmor turns. Another novelty in treating the resonance and relaxation problems in a ferromagnet nanoparticle assembly is that we consider the dipole interactions between magnetic moments in a microscopic way, not by introducing any phenomenological relaxation time.

ACKNOWLEDGMENT

Work was done under the auspices of The Russian Foundation for Basic Research projects 10-02-96023 and 11-07-96007.

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