Abstract: A spin system is considered with a Hamiltonian typical of molecular magnets, having dipole-dipole interactions and a single-site magnetic anisotropy. In addition, spin interactions through the common radiation field are included. A fully quantum-mechanical derivation of the collective radiation rate is presented. An effective narrowing of the dipole-dipole attenuation, due to high spin polarization is taken into account. The influence of the radiation rate on spin dynamics is carefully analysed. It is shown that this influence is completely negligible. No noticeable collective effects, such as superradiance, can appear in molecular magnets, being caused by electromagnetic spin radiation. Spin superradiance can arise in molecular magnets only when these are coupled to a resonant electric circuit, as has been suggested earlier by one of the authors in Laser Phys. 12, 1089 (2002).

Absence of spin superradiance in resonatorless magnets

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1. Introduction

Superradiance is a phenomenon well known for atomic systems [1,2], where it occurs because of the self-organized correlation of atomic transitions through the common radiation field. There exists a similar phenomenon of spin superradiance arising in spin systems, as is reviewed in [3,4]. An accurate microscopic theory of spin superradiance has recently been developed [5–11] demonstrating that, despite many similarities between the atomic and spin superradiance, these phenomena possess features that radically distinguish the spin from atomic superradiance [3–11]. The main such a fundamental difference is that the spin superradiance in real systems cannot be caused by their magnetodipole radiation field. The fact that the spin radiation rate is negligible as compared to all other relaxation rates was, first, noticed by Bloembergen [12]. This radiation rate cannot yield the development of spin superradiance, as has been emphasized recently [13] and analysed in detail in review [4]. In order to realize spin superradiance, it is necessary to couple a polarized magnet to a resonant electric circuit. The collectivization of spin motion happens due to the resonator feedback field, which, thus, replaces the common radiation field [3–11]. The analogies and differences in atomic and spin superradiance have been extensively analysed in reviews [3,4,14].

Even though the physics of spin superradiance has been well established, there have recently appeared discussions on the role of the radiation rate in spin relaxation. The interest to this problem has been resumed in
connection to the studies of a novel class of materials, called molecular magnets, which can be highly polarized and may possess large molecular spins (see review [4]). The possibility of realizing spin superradiance by molecular magnets, coupled to a resonator, was advanced in Ref. [10], being based on the well developed theory [3–11]. At the same time, there have appeared speculations that spin superradiance in molecular magnets could be produced without coupling them to a resonant circuit, but in a way completely equivalent to atomic superradiance, so that the collectivization in spin motion would be due solely to their interactions through the common radiation field. Plausible collective effects in molecular magnets, caused by the radiation interaction, were discussed in Ref. [15], though the consideration there involved a very strong transverse magnetic field. At this instance, it is necessary to stress that superradiance as such is defined as a self-organized process, without being pushed by strong transverse fields. This definition was given in the original paper by Dicke [16] and is generally accepted for both atomic [1,2,17] as well as spin systems [3,4]. Having a generally accepted definition, there is no reason to change it. Therefore a process, involving strong transverse fields, has nothing to do with superradiance. In the best case, this could be collective spin induction, provided there would really be any collective effects. The possibility of the collective spin induction, caused by the molecular interactions through the electromagnetic radiation they are emitting, was considered in Ref. [15] in the frame of the phenomenological Landau-Lifshits-type equation for the classical magnetic moment, without taking account of the dipole spin interactions. However these dipole interactions do exist and are very strong in all really available molecular magnets.

An effort of taking into account the dipole interactions was attempted in Ref. [18]. Unfortunately, the authors confused the effective spin temperature with the real temperature of the lattice. These temperatures, as is well known [19–21], have nothing in common, when the system is not in equilibrium. The lattice temperature is intrinsically positive, while the spin temperature is positive or negative depending upon the average expectation value of spin polarization. Moreover, there even exist two effective spin temperatures for the same spin system, the Zeeman and dipole effective temperatures. These temperatures are directly related to the average spin, which in a nonequilibrium system is a function of time. As a consequence, the spin temperatures are also functions of time and are defined by the Provotorov evolution equations. For inverted nonequilibrium spins, the Zeeman temperature is negative. One should not confuse the negative-time-dependent effective Zeeman spin temperature with the positive stationary real lattice temperature.

In the present paper, we give an accurate analysis of the dynamics for a spin system described by a microscopic Hamiltonian typical of molecular magnets, including the dipole interactions together with the electromagnetic-field interaction. All consideration is based on a fully quantum-mechanical picture and on well grounded mathematical methods of treating nonlinear spin-evolution equations. We do not couple the magnet to a resonant circuit, which allows us to concentrate our attention on the role of the radiation interactions and on the influence of the related radiation rate on spin motion. Our firm conclusion is that these radiation interactions play no role and the radiation rate is negligible, but the spin dynamics is governed by the dipole spin interactions. In this way, the spin interaction through the radiation field is absolutely unable to produce superradiance in molecular magnets.

2. Electromagnetic spin interactions

A realistic microscopic Hamiltonian for molecular magnets has the form

\[
\hat{H} = \sum_{i=1}^{N} \hat{H}_i + \frac{1}{2} \sum_{i<j}^{N} \hat{H}_{ij} + \hat{H}_f. \tag{1}
\]

Here the first term describes the single-spin energy, with

\[
\hat{H}_i = -\mu_0 \mathbf{S}_i \cdot \mathbf{B}_i - D (\mathbf{S}_i^2)^2, \tag{2}
\]

where \(\mu_0 = \hbar \gamma_S \gamma_S\) being the gyromagnetic ratio for a molecule of spin \(S\); \(D\) is a single-site anisotropy parameter; the total magnetic field

\[
\mathbf{B}_i = B_0 \mathbf{e}_z + \mathbf{H}_i, \tag{3}
\]

acting on each spin, contains an external magnetic field \(B_0\) and the radiation field

\[
\mathbf{H}_i \equiv \mathbf{H}(\mathbf{r}_i, t). \tag{4}
\]

The second term in Eq. (1) corresponds to dipole spin interactions, with

\[
\hat{H}_{ij} = \sum_{\alpha\beta} D_{ij}^{\alpha\beta} \mathbf{S}_i^\alpha \mathbf{S}_j^\beta, \tag{5}
\]

where

\[
D_{ij}^{\alpha\beta} \equiv \frac{\mu_0^2}{r_{ij}^3} \left( \delta_{\alpha\beta} - 3n_{ij}^\alpha n_{ij}^\beta \right)
\]

is the dipolar tensor and

\[
r_{ij} \equiv |\mathbf{r}_{ij}|, \quad n_{ij}^\alpha \equiv \frac{\mathbf{r}_{ij}^\alpha}{r_{ij}}, \quad r_{ij} \equiv \mathbf{r}_i - \mathbf{r}_j.
\]

The last term of Eq. (1) is the Hamiltonian of electromagnetic field

\[
\hat{H}_f = \frac{1}{8\pi} \int \left( \mathbf{E}^2 + \mathbf{H}^2 \right) \, d\mathbf{r}, \tag{6}
\]

with an electric field \(\mathbf{E} = \mathbf{E}(\mathbf{r}, t)\) and magnetic radiation field \(\mathbf{H} = \mathbf{H}(\mathbf{r}, t)\). The latter is expressed through the vector potential \(\mathbf{A}\) as \(\mathbf{H} = \nabla \times \mathbf{A}\). In what follows, the Coulomb calibration \(\nabla \cdot \mathbf{A} = 0\) is used.
To write the equations of motion for the spin operators $S_i$, in a compact form, it is convenient to introduce the notation for the local fluctuating fields

$$\xi \equiv \frac{1}{\hbar} \sum_{j(\neq i)} (a_{ij} S_j^- + c_{ij} S_j^+ + c_{ij} S_j^z),$$

$$\xi \equiv \frac{i}{\hbar} \sum_{j(\neq i)} (2c_{ij} S_j^z - \frac{1}{2} a_{ij} S_j^x + 2b_{ij} S_j^y),$$

in which $S_j^\pm$ are the ladder spin operators and

$$a_{ij} \equiv D_{ij}^z, \quad b_{ij} \equiv \frac{1}{4} (D_{ij}^{xx} - D_{ij}^{yy} - 2iD_{ij}^{yz}),$$

$$c_{ij} \equiv \frac{1}{2} (D_{ij}^{xx} - iD_{ij}^{yz}).$$

The magnetic moment operator can be represented as

$$\mathbf{M}_t \equiv \mu_0 \mathbf{S}_t = \mu S_t^x + \mu^* S_t^+ + \mu_0 S_t^z,$$

with the employed notation

$$\mu \equiv \frac{\mu_0}{2} (\mathbf{e}_x - i\mathbf{e}_y), \quad \mu_0 \equiv \mu_0 \mathbf{e}_z.$$

The latter vectors enjoy the properties

$$|\mu|^2 = \frac{\mu_0^2}{2}, \quad \mu^2 = 0, \quad \mu \cdot \mu_0 = 0.$$

And two more important notations are the Zeeman frequency

$$\omega_0 \equiv -\frac{1}{\hbar} \mu_0 B_0$$

and the effective field

$$f \equiv -i 2\mu \cdot \mathbf{H}_t + \xi.$$

Then the equations of motion for the spin operators read as

$$\frac{dS_i^-}{dt} = -i \left[ \omega_0 - \frac{1}{\hbar} \mu_0 \cdot \mathbf{H}_t + \xi_0 \right] S_i^- +$$

$$+ f S_i^+ + \frac{i}{\hbar} D \left( S_i^- S_i^+ + S_i^z S_i^- \right),$$

$$\frac{dS_i^+}{dt} = \frac{1}{2} \left( f^+ S_i^- + S_i^+ f \right).$$

The equations for the electromagnetic field are

$$- \frac{1}{c} \frac{\partial \mathbf{E}}{\partial t} = \nabla \times \mathbf{H} - \frac{4\pi}{c} j, \quad - \frac{1}{c} \frac{\partial \mathbf{A}}{\partial t} = -\mathbf{E},$$

where the current density is

$$j = -c^2 \sum_{i=1}^{N} \mathbf{M}_t \times \nabla \delta(\mathbf{r} - \mathbf{r}').$$

From Eqs. (11) one has the equation

$$\left( \nabla^2 - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \right) \mathbf{A} = -\frac{4\pi}{c} j,$$

whose solution is

$$\mathbf{A}(\mathbf{r}, t) = \frac{1}{c} \int \mathbf{j} \left( r', t - |\mathbf{r} - \mathbf{r}'| \right) \frac{d\mathbf{r}'}{|\mathbf{r} - \mathbf{r}'|}.$$

The average vacuum fluctuations are assumed to be zero. The vector potential (14), with the current density (12), can be represented as the sum

$$\mathbf{A}_i = \mathbf{A}_i^- + \mathbf{A}_i^+ + \mathbf{A}_i',$$

in which

$$\mathbf{A}_i^- = -\sum_j \left[ 1 + \frac{i}{c} \frac{\partial}{\partial t} \right] \frac{r_{ij}}{r_{ij}^2} \times \mu^* S_j^- \left( t - \frac{r_{ij}}{c} \right),$$

$$\mathbf{A}_i' = -\sum_j \frac{r_{ij}}{r_{ij}^3} \times \mu_0 S_j^z \left( t - \frac{r_{ij}}{c} \right),$$

where the property

$$\frac{\partial}{\partial t} S_i^- \left( t - \frac{r_{ij}}{c} \right) = -\frac{1}{c} \frac{\partial}{\partial t} S_i^- \left( t - \frac{r_{ij}}{c} \right)$$

is used. From Eqs. (15) and (16), one gets

$$\mathbf{H}_i = \nabla \times \mathbf{A}_i = \mathbf{H}_i^- + \mathbf{H}_i^+ + \mathbf{H}_i',$$

with the expressions

$$\mathbf{H}_i^- = -\sum_j \left[ \mu^* - (\mu^* \cdot \mathbf{n}_{ij}) \mathbf{n}_{ij} \frac{\partial^2}{\partial t^2} + \frac{\mu^* - 3(\mu^* \cdot \mathbf{n}_{ij}) \mathbf{n}_{ij}}{r_{ij}^3} \left( 1 + \frac{r_{ij}}{c} \frac{\partial}{\partial t} \right) \right] S_j^- \left( t - \frac{r_{ij}}{c} \right),$$

$$\mathbf{H}_i' = -\sum_j \frac{\mu_0 - 3(\mu_0 \cdot \mathbf{n}_{ij}) \mathbf{n}_{ij}}{r_{ij}^3} S_j^z \left( t - \frac{r_{ij}}{c} \right).$$

Keeping in mind a macroscopic in all directions sample, we may simplify Eqs. (18) by averaging them over the spherical angle $\Omega(\mathbf{n}_{ij})$ related to the unit vector $\mathbf{n}_{ij}$. In doing this, we employ the integrals

$$\frac{1}{4\pi} \int \mu - (\mu \cdot \mathbf{n}) \mathbf{n} d\Omega(\mathbf{n}) = \frac{2}{3} \mu,$$

$$\frac{1}{4\pi} \int \mu - 3(\mu \cdot \mathbf{n}) \mathbf{n} d\Omega(\mathbf{n}) = 0,$$

$$\frac{1}{4\pi} \int (\mu \cdot \mathbf{n}) \mathbf{n} d\Omega(\mathbf{n}) = \frac{1}{3} \mu.$$

Then from Eqs. (18), we find

$$\mathbf{H}_i^- = -\frac{2}{3} \mu^* \sum_j \frac{1}{c^2 r_{ij}^3} \frac{\partial^2}{\partial t^2} S_j^- \left( t - \frac{r_{ij}}{c} \right).$$
and $H'_i = 0$.

Since we are looking for the possibility of arising coherent collective effects, we have to accept the first basic condition for the existence of such effects. This necessary condition is the assumption that the transverse spin motion can be characterized by a well-defined frequency $\omega$, with the related wavelength $\lambda \equiv 2\pi/k$ and wave vector $k = \omega/c$. Under this condition, the Born approximation is valid,

$$S'_j \left( t - \frac{r}{c} \right) = S_j(t) \Theta(ct - r)e^{ikr}, \quad (20)$$

where $\Theta(\cdot)$ is a unit-step function, which allows us to rewrite Eq. (19) in the form

$$H'_i = \frac{2}{3} k^3 \mu^* \sum \exp(ikr_{ij}) \Theta(ct - r_{ij}) S'_j(t). \quad (21)$$

The magnetic field (17), with taking account of Eq. (21) and of the properties $\mu \cdot H'_i = 0$ and $\mu_0 H'_i = 0$, has to be substituted to the effective force (9) entering the spin equations of motion (10). The latter equations can be treated by the scale separation approach [5–10,14,22], whose mathematical foundation is based on the averaging techniques [23]. In this approach, one treats the local fluctuating fields (6) as random variables, which allows for the account of quantum spin fluctuations due to dipole interactions. Then the equations of motion (10) are to be averaged over the spin degrees of freedom. To this end, we define the averages for the transverse spin

$$u(r, t) \equiv \frac{1}{S} \langle S^-(r, t) \rangle, \quad (22)$$

coherence intensity

$$w(r, t) \equiv \frac{1}{S^2} \langle S^+(r, t)S^-(r + 0, t) \rangle, \quad (23)$$

and spin polarization

$$s(r, t) \equiv \frac{1}{S} \langle S^z(r, t) \rangle. \quad (24)$$

To proceed further, one needs to invoke the second fundamental condition for the occurrence of collective effects. This is the existence of the regions of strongly correlated spins, which are called spin packets. Really, without a substantial correlation, no collective effects can arise. The existence of a strongly correlated spin motion in some regions of the sample implies that the spin motion inside each of the regions is coherent, while the correlation between different regions of correlation is either absent or very weak. Let the characteristic length of a correlated region be $L_c$, volume $V_c = L^3_c$, and the number of spins inside it $N_c$. This characteristic length has to satisfy the inequality

$$kL_c \ll 1. \quad (25)$$

If inequality (25) does not hold, then all spins in sum (21) oscillate independently of each other, with their correlation limited by only nearest neighbours, which follows from the fast spatial oscillations of the kernel in sum (21). But under condition (25), this kernel is practically constant, and all spins in the region $V_c$ can be synchronized. When the wavelength $\lambda \gg L$ is much larger than the system length, then $L_c = L$, and the whole system radiates coherently. This, however, is not compulsory. The wavelength $\lambda$ can be much smaller than $L$. In that case, the system separates into several correlated regions, which radiate almost independently from each other. As a result, the total radiation pulse has an oscillatory behaviour. This is a well known picture in optical coherent radiation [1,2]. Lasers with a large aperture always radiate not by a unique beam but by a bunch of filaments [24–28]. The same is true for spin systems. When $\lambda \ll L$, the sample is divided into correlated regions of size $L_c \ll L$, but such that $kL_c \ll 1$.

Inside of a correlated region, the average ($S'_j$) weakly depends on the site $j$, so that this average can be taken out of the sum. Thus, we come to the definition of the collective radiation rate

$$\gamma_r \equiv \gamma_0 \sum_{j} N_c \sin(kr_{ij}) \Theta(ct - r_{ij}) \quad (26)$$

and of the collective frequency shift

$$\delta \omega \equiv \gamma_0 \sum_{j} N_c \cos(kr_{ij}) \Theta(ct - r_{ij}), \quad (27)$$

in which

$$\gamma_0 \equiv \frac{2}{3h} \mu_0^2 Sk^3 \quad (28)$$

is the single-spin natural width. Inside such a correlated region, one has

$$-i \frac{\hbar}{2}\langle \mu \cdot H_i \rangle = (\gamma_r - i\delta \omega)\alpha.$$

Averaging the spin equations of motion (10) over the spin degrees of freedom, we decouple the binary spin expressions $\langle S^x_i S^y_j \rangle$, for $i \neq j$, as $\langle S^x_i \rangle \langle S^y_j \rangle$, treating the local fluctuating fields (6) as random variables. At the same time, the spin products for the coinciding sites must be separated into several correlated regions, which radiate almost independently from each other. The equations of motion (26) and (27), manifest themselves at the time scale $L_c/c$. Taking for $L_c$ the maximal value $L_c \sim 1/k$,
the retardation time can be approximated by $\omega^{-1}$. Thus, Eqs. (26) and (27) can be represented as

$$\gamma_r \cong \gamma_0 N_e \left(1 - e^{-\omega^{-1}}\right),$$

$$\delta \omega \cong \gamma_0 \frac{3N_e}{2kL_c} \left(1 - e^{-\omega^{-1}}\right).$$

The time $1/\omega$ is very short, usually it is the shortest among all other relaxation times. Therefore, for $\omega t \gg 1$, and taking into account that $N_e = \rho L_c^3$, we have for the collective radiation rate

$$\gamma_r \cong \gamma_0 N_e \frac{2}{3\hbar} \rho_0^2 S k^3 N_c,$$

and for collective frequency shift

$$\delta \omega \cong \frac{3\gamma_r}{2kL_c} = \frac{1}{\hbar} \rho_0^2 S (kL_c)^2,$$

where $\rho \equiv N/V = N_c/V_c$ is spin density. Expressions (30) and (31) were, first, found by Ginzburg [30], who used a classical picture, which also was discussed in [4,31]. Here we have presented a fully quantum-mechanical derivation of these expressions and, in addition, have obtained their generalizations (26), (27), and (29).

In studying spin dynamics, it is necessary to take account of the longitudinal and transverse relaxation rates. The former is due to spin-lattice interactions, and is denoted as $\gamma_1$. The transverse relaxation rate is caused by the presence of dipole spin interactions and its standard value is

$$\gamma_2 = \frac{n_0}{\hbar} \rho_0^2 S \sqrt{S(S+1)},$$

where $n_0$ is the number of nearest neighbours. Quantity (32) is derived [19–21] under the assumption of a small longitudinal spin polarization. When the latter is large, so that the average polarization (24) substantially differs from zero, then the effective transverse relaxation rate narrows, becoming dependent on $s$. This effective narrowed rate was calculated and described in detail by Abragam and Goldman [21], who express its value

$$\gamma_2(s) = \frac{CM_2(s)}{M_4(s) - M_2^2(s)}^{1/2}$$

through the moments

$$M_2(s) = M_2(0)(1 - s^2),$$

$$M_4(s) = 2.18M_2^2(s)(1 - s^2)(1 - 0.42s^2).$$

The constant $C$ in Eq. (33) depends on the line shape, being $C = \pi/2$ for the Gaussian and $C = \pi^2/2$ for Lorentzian lines. Substituting these moments in Eq. (33) and keeping in mind that $s^2 \leq 1$, we obtain

$$\gamma_2(s) = \gamma_2(0)(1 - s^2),$$

where $\gamma_2(0) = \sqrt{CM_2(0)} \equiv \gamma_2$. The total transverse relaxation rate, including the inhomogeneous broadening $\gamma_2$, is

$$\Gamma_2 = \gamma_2(s) + \gamma_2^*.$$

Averaging Eqs. (10) over the spin degrees of freedom, we take into account all described attenuations. Then, using the notation for the anisotropy frequency

$$\omega_D \equiv \frac{1}{\hbar} (2S - 1)D,$$

and for the effective spin frequency

$$\omega_s \equiv \omega_0 - \omega_D s,$$

we come to the evolution equations

$$\frac{du}{dt} = -i(\omega_s + \xi_0 - i\Gamma_2) u + fs,$$

$$\frac{dw}{dt} = -2\Gamma_2 w + (u^* f + f^* u)s,$$

$$\frac{ds}{dt} = -\frac{1}{2} (u^* f + f^* u) - \gamma_1 (s - \xi),$$

in which the effective force is

$$f = (\gamma_r - i\delta \omega) u + \xi.$$

Because of the occurrence of the random variables $\xi_0$ and $\xi$, corresponding to local spin fluctuations, these are stochastic differential equations. Averaging over the random fluctuations, treated as a Gaussian white noise [32], with the width $\gamma_1$, we follow the scale separation approach [5–10,14,22]. In this way, we obtain the equations for the guiding centers

$$\frac{dw}{dt} = -2(\gamma_2 + \gamma_2^* - \gamma_2 s^2 - \gamma_1 s)w + 2\gamma_3 s^2,$$

$$\frac{ds}{dt} = -\gamma_1 w - \gamma_3 s - \gamma_1 (s - \xi),$$

containing all relaxation rates discussed above.

Comparing the relaxation rates (30) and (32), we have

$$\frac{\gamma_2}{\gamma_2} = \frac{2}{3n_0} \sqrt{S/S + 1} (kL_c)^3.$$

According to condition (25) for the existence of correlated regions, $kL_c \ll 1$, and since $n_0 \sim 10$, the radiation rate $\gamma_r$ is always much smaller than the dipole relaxation rate $\gamma_2$. As is clear from Eqs. (39), the value $\gamma_r/\gamma_2$ plays no role in spin dynamics. Even taking the maximal correlation length $L_c \sim 1/k$, one has $\gamma_r/\gamma_2 \sim 0.1$, which does not noticeably influence the spin relaxation. Equations (39) contain no solutions corresponding to spin superradiance.
The longitudinal relaxation rate was assumed to be small, \( \gamma_1 \ll \gamma_2 \), which is usually the case. Other relaxation rates, expressed in units of \( \gamma_2 \), were varied in a wide diapason: \( \gamma_2^* \in [0,1] \), \( \gamma_3 \in [0,1] \), and \( \gamma_r \in [0,1] \). Different initial conditions \( w_0 = w(0) \) and \( s_0 = s(0) \) were considered, satisfying the inequality \( w_0 + s_0^2 \leq 1 \). The solutions for \( \gamma_r = 0 \) and \( \gamma_r = 0.1 \) were found to be practically indistinguishable. Even unrealistically large \( \gamma_r = 1 \) resulted in insignificant changes, as compared to solutions with \( \gamma_r = 0 \). The main rates, governing the spin evolution, are \( \gamma_2 \), \( \gamma_2^* \), and \( \gamma_3 \). In Figs. 1–4, we show the solutions for \( w(t) \) and \( s(t) \) at different relaxation rates, demonstrating the role of the latter, for the same initial conditions \( w_0 = 0 \) and \( s_0 = 1 \). These figures display the effect of self-induced dynamical coherence caused by the existence of dipole interactions. The main term in Eq. (39), producing the coherence intensity, is that containing \( \gamma_3 \). Without this term, no maximum in \( w \) appears. The term with \( \gamma_2 s_0^2 \), due to the narrowing of the effective transverse relaxation, which results from a high spin polarization, is much less important and influences the spin motion by about 10%. If at the initial moment of time, an external transverse field imposes a noticeable coherence intensity \( w_0 > 0 \), then we have the standard spin induction, as shown in Fig. 5.

Both the spin induction and self-induced dynamical coherence occur on the time scale of the order of \( T_2 \). Therefore, though these are coherent phenomena, they have nothing to do with superradiance which requires the pulse time to be much shorter than \( T_2 \). Also, there is no spin reversal typical of superradiance.

Recently, there have been attempts to detect the electromagnetic radiation generated in the avalanches of mag-
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