Nonlinear and thermal effects in the absorption of microwaves by random magnets

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We study the temperature dependence of the absorption of microwaves by random-anisotropy magnets. It is governed by strong metastability due to the broad distribution of energy barriers separating different spin configurations. At a low microwave power, when the heating is negligible, the spin dynamics is close to linear. It corresponds to the precession of ferromagnetically ordered regions that are in resonance with the microwave field. Previously, we showed [Phys. Rev. B **103**, 214414 (2021)] that in this regime a dielectric substance packed with random magnets would be a strong microwave absorber in a broad frequency range. Here we demonstrate that on increasing the power, heating and over-barrier spin transitions come into play, resulting in the nonlinear behavior. At elevated temperatures the absorption of microwave power decreases dramatically, making the dielectric substance with random magnets transparent for the microwaves.

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I. INTRODUCTION

Random-anisotropy (RA) magnets have numerous technological applications. Their static properties have been intensively studied during the last four decades (see, e.g, Refs. [1–3] and references therein). They are materials that have ferromagnetic exchange interaction between neighboring spins but random directions of local magnetic anisotropy, either due to their amorphous structure or due to sintering from randomly oriented nanocrystals. Depending on the ratio of the ferromagnetic exchange J and local magnetic anisotropy D_R, they can be hard or soft magnets. At $D_R \ll J$ they consist of large ferromagnetically oriented regions, often called Imry-Ma (IM) domains, and are characterized by low coercivity and high magnetic susceptibility. In the opposite limit of large anisotropy compared to the exchange they have high coercivity and a large area of the hysteresis loop. The second limit of $D_R \gtrsim J$ would be difficult to achieve in conventional ferromagnets due to the spin-orbit relativistic nature of the magnetic anisotropy. However, in a sintered RA magnet it is the effective magnetic anisotropy $\tilde{D}_R$ of the nanocrystal that enters the problem. Since $\tilde{D}_R$ goes up with the average size of the nanocrystal (or an amorphous structure factor), both limits can easily be realized in experiments.

The case of strong effective magnetic anisotropy is conceptually simple. It is equivalent to the array of densely packed, randomly oriented, single-domain magnetic particles. However, the case of $D_R \lesssim J$ has been the subject of a significant controversy. It was initially analyzed in terms of the IM argument [4,5] that explores the analogy with the random walk problem: Weak random local pushes from the RA make the direction of the magnetization created by the strong exchange interaction wander around the magnet with the ferromagnetic correlation length $R_f/a \propto (J/D_R)^{2/(4-d)}$ (a is the interatomic distance, and d is the dimensionality of the system). The validity of this argument that ignores metastable states was later questioned by numerical studies [6–8]. It was found that RA magnets exhibit metastability and history dependence [9,10] regardless of the strength of the RA, although the IM argument roughly holds for the average size of ferromagnetically correlated regions if one begins with a fully disordered state. More recently, using the random-field model, it was demonstrated that the presence or absence of topological defects determined by the relation between the number of spin components and dimensionality of space is crucial for the properties of random magnets [11–13]. Nevertheless, questions about the exact ground state, spin-spin correlations, topological defects, etc., in the RA model remain largely unanswered after a 40-year effort.

In the absence of a rigorous theory describing static properties of the RA ferromagnet, studies of the dynamics present a significant challenge. Collective modes and their localization have been observed in amorphous ferromagnets with random local magnetic anisotropy [14–18], inhomogeneous thin magnetic films [19], and submicron magnetic heterostructures [20] and in films where an inhomogeneous magnetic field was generated by the tip of a force microscope [21]. The complex nature of these excitations, which involves longitudinal, transversal, and mixed modes, has been reported. Following these experiments, an analytical theory of the uniform spin resonance in a thin film of the RA ferromagnet in a nearly saturating magnetic field was developed [22]. The dependence of the frequency of the longitudinal resonance on the magnetic field and its angle with the film were obtained and compared with experimental findings.

Practical applications of random magnets as absorbers of microwave power typically involve very weak static magnetic fields. This problem is more challenging than the problem with the saturating field because the underlying magnetic state is strongly disordered, resembling a spin glass. Zero-field resonances were observed in spin glasses in the past [23–27]. They were attributed [28–30] to the RA arising from

Dzyaloshinskii-Moriya interaction and analyzed within hydrodynamic theory [31,32]. Due to the lack of progress on spin glasses and random magnets, theoretical effort aimed at understanding their dynamics was largely abandoned in the 1990s. Recently, we returned to this problem using the power of modern computers within the framework of the RA magnet in zero magnetic field [33]. Computer-generated images of spin oscillations induced by the ac field were obtained. They confirmed earlier conjectures that collective modes in the RA ferromagnet were localized within correlated volumes that are in resonance with the ac field. It was shown that broad distribution of resonances makes such magnets excellent broadband absorbers of the microwave power that can compete with nanocomposites commonly used for that purpose [34,35].

In Ref. [33] the power of the microwave radiation was assumed to be too weak to cause any heating or significant nonlinearity in the response of a system consisting, e.g., of a dielectric layer packed with RA magnets. Meanwhile, in certain applications such a system would be subjected to a pulse of a strong microwave beam from a nearby source that might cause nonlinear effects and heating. The problem studied in this paper, besides its fundamental interest, is related to the practical question of how a short, powerful microwave signal would modify the absorption of a continuous weak microwave signal coming from an independent distant source.

In the RA magnet, the absorbed energy is quickly redistributed over all degrees of freedom, so that the magnetic system is effectively in equilibrium. Consequently, the nonlinearity and saturation of the absorption can be effectively described in terms of heating. Then the task effectively reduces to the computation of the temperature dependence of the power absorption. This allows one to describe the whole process in terms of a single differential equation for the time dependence of the spin temperature. We use the Monte Carlo method to prepare initial states at different temperatures, then run a conservative dynamics and calculate the temperature-dependent absorbed power using the fluctuation-dissipation theorem (FDT). The frequency dependence of the power shows the evolution from a broad absorption peak at low temperatures to the low-absorption plateau on increasing the temperature.

Practically, there is energy transfer from the spins to the atomic lattice and further to the dielectric matrix containing the RA magnets and to the substrate that it is deposited on. The flow of energy from spins to other degrees of freedom they are interacting with inside the magnet, such as phonons, is fast and effectively increases the heat capacity of the system, reducing the heating to some extent. The heat flow from the magnets to the dielectric matrix and then to the substrate is much slower. Here we consider high-power microwave pulses of a duration that is sufficient to heat the RA magnet but short compared to the typical times of heat flow out of the RA magnet due to thermal conductivity. We show that such a pulse can make the RA magnets transparent for the microwaves.

This paper is organized as follows. The RA model and relevant theoretical concepts are introduced in Sec. II. Section III spells out two numerical experiments performed in this work: (i) pumping the system by an ac field and (ii) running conservative dynamics at different temperatures and computing the absorbed power by the FDT. Section IV reports numerical data from these two experiments. Our results and possible applications are discussed in Sec. V.

II. THEORY

Following Ref. [33], we consider a model of three-component classical spin vectors $\mathbf{s}_i$ on the lattice described by the Hamiltonian

$$\mathcal{H} = -\frac{1}{2}\sum_{i,j} J_{ij}\mathbf{s}_i \cdot \mathbf{s}_j - \frac{D_R}{2}\sum_i (\mathbf{n}_i \cdot \mathbf{s}_i)^2 - \mathbf{h}(t)\cdot\sum_i \mathbf{s}_i. \quad (1)$$

Here the first term is the exchange interaction between nearest neighbors with the coupling constant $J > 0$, D_R is the strength of the easy-axis RA in energy units, $\mathbf{n}_i$ is a three-component unit vector having random direction at each lattice site, and $\mathbf{h}(t) = \mathbf{h}_0 \sin(\omega t)$ is the ac magnetic field in energy units. We study situations in which the wavelength of the electromagnetic radiation is large compared to the size of the RA magnet, so that the time-dependent field acting on the spins is uniform across the system. In all practical cases, h_0 is small in comparison to the other terms of the Hamiltonian, so that nonlinearity can emerge from only the saturation of resonances.

We assume [1,5,12] that the behavior of the system is dominated by the exchange and random anisotropy and neglect the dipole-dipole interaction (DDI) between the spins, which is typically weaker. Consequently, the IM domains generated by the RA are much smaller than typical magnetic domains generated by the DDI [2]. The study of RA magnets requires systems with size greater than the ferromagnetic correlation length. Adding DDI to such a problem considerably slows down the numerical procedure without changing the results in any significant way. According to Ref. [33], the microwave absorption is qualitatively similar in RA systems of different dimensions. With that in mind and to reduce computation time that becomes prohibitively long for a large three-dimensional (3D) system, we do most of the numerical work for the two-dimensional (2D) model on a square lattice.

We consider conservative dynamics of the magnetic system governed by the dissipationless Landau-Lifshitz equation

$$\hbar\dot{\mathbf{s}}_i = \mathbf{s}_i \times \mathbf{h}_{\mathrm{eff},i}, \quad \mathbf{h}_{\mathrm{eff},i} \equiv -\frac{\partial\mathcal{H}}{\partial\mathbf{s}_i} = \mathbf{h}(t) + \mathbf{H}_{\mathrm{eff},i}, \quad (2)$$

where $\mathbf{H}_{\mathrm{eff},i} = \sum_j J_{ij}\mathbf{s}_j + D_R(\mathbf{n}_i\cdot\mathbf{s}_i)\mathbf{n}_i$ is the effective field due to the exchange and RA. It was shown in Ref. [33] that including a small dissipation due to the interaction of the magnetic system with other degrees of freedom does not change the results for the absorbed power. This occurs due to the fact that the RA magnet has a continuous spectrum of resonances and the system has its own internal damping due to the many-particle nature of the Hamiltonian, Eq. (1), and the distribution of energy between different modes. For the conservative system, the time derivative of $\mathcal{H}$ is equal to the absorbed power:

$$\dot{\mathcal{H}}(t) = P_{\mathrm{abs}} = -\dot{\mathbf{h}}(t)\cdot\sum_i \mathbf{s}_i(t). \quad (3)$$

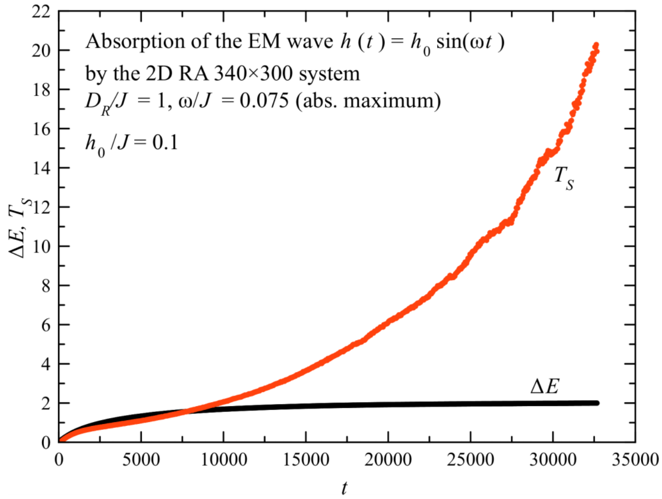

FIG. 1. Time dependence of the absorbed microwave energy ΔE and of the spin temperature T_s in the RA ferromagnet.

An important component of the theory of classical spin systems is the dynamical spin temperature [36]. For the RA magnet the general expression given in Ref. [36] yields [37]

$$T_S = \frac{\sum_i (\mathbf{s}_i \times \mathbf{H}_{\mathrm{eff},i})^2}{2\sum_{i,j} J_{ij}\mathbf{s}_i \cdot \mathbf{s}_j + D_R \sum_i [3(\mathbf{n}_i \cdot \mathbf{s}_i)^2 - 1]}. \tag{4}$$

At $T = 0$, spins are aligned with their effective fields, and the numerator of this formula vanishes. At $T = \infty$, spins are completely disordered, and both terms in the denominator vanish. These two limiting cases are clearly seen in Fig. 1. In a thermal state of a large system with temperature T, created by Monte Carlo, $T_S \cong T$ up to the small fluctuations.

At low temperatures and weak RA, $D_R \ll J$, spins are strongly correlated within large regions of the characteristic size R_f (ferromagnetic correlation radius) defined by [4,5]

$$\frac{R_f}{a} \sim \left(\frac{J}{D_R}\right)^{2/(4-d)}, \tag{5}$$

where a is the lattice spacing. In three dimensions, R_f becomes especially large, which makes computations for systems with linear size $L \gg R_f$ problematic. Assuming that there is no long-range order (which is the case when the magnetic state is obtained by energy minimization from random initial directions of spins), one can estimate R_f using the value of the average spin

$$\mathbf{m} = \frac{1}{N}\sum_i \mathbf{s}_i, \tag{6}$$

which is nonzero in finite-size systems due to fluctuations. One has

$$m^2 = \frac{1}{N^2}\sum_{i,j} \mathbf{s}_i \cdot \mathbf{s}_j = \frac{1}{N}\sum_j \langle \mathbf{s}_i \cdot \mathbf{s}_{i+j}\rangle \Rightarrow \frac{1}{N}\int_0^\infty \frac{d^d r}{a^d} G(r), \tag{7}$$

where $G(r)$ is the spatial correlation function and d is the dimensionality of the space. As the RA magnet has lots of metastable local energy minima, $G(r)$ depends on the initial conditions and on the details of the energy minimization routine. In two dimensions for $G(r) = \exp[-(r/R_f)^p]$ one obtains

$$m^2 = K_p \frac{\pi R_f^2}{N a^2} \quad \Longrightarrow \quad \frac{R_f}{a} = m\sqrt{\frac{N}{\pi K_p}}, \tag{8}$$

where $K_1 = 2$ and $K_2 = 1$.

Having estimated R_f, one can find the number of IM domains N_{IM} in the system of size used in the numerical work. In 2D with linear sizes L_x and L_y one has $N_{\mathrm{IM}} = L_x L_y/(\pi R_f^2)$. In particular, for a system with $N = 300 \times 340 = 102\,000$ spins and $D_R/J = 0.3$, energy minimization at $T = 0$ starting from a random spin state yields $m \approx 0.21$, and with $p = 2$ one obtains $R_f/a \approx 37.8$ and $N_{\mathrm{IM}} \approx 23$. For $D_R/J = 1$, one obtains $m \approx 0.074$ and $R_f/a \approx 13.3$, which yields $N_{\mathrm{IM}} \approx 183$. For the ratio of the R_f values one obtains $R_f^{(D_R=0.3)}/R_f^{(D_R=1)} \approx 2.84$, which is close to the value of 3.33 given by Eq. (5).

In the linear regime the ac field generates a deviation of the magnetization from its initial value,

$$m_x(t) = h_0[\chi_x'(\omega)\sin(\omega t) - \chi_x''(\omega)\cos(\omega t)]. \tag{9}$$

Here $\mathbf{h}_0 = h_0\mathbf{e}_x$, and χ is the susceptibility. At finite temperatures the absorbed power in this regime can be computed using the fluctuation-dissipation theorem (FDT), which relates the imaginary part of the susceptibility to the Fourier transform of the autocorrelation function of the average spin:

$$\chi_x''(\omega) = \frac{\omega N}{k_B T}\mathrm{Re}\int_0^\infty dt\, e^{i\omega t} A_x(t), \tag{10}$$

where

$$A_x(t) \equiv \langle [m_x(t) - \langle m_x\rangle][m_x(0) - \langle m_x\rangle]\rangle. \tag{11}$$

The absorbed power per spin is related to $\chi_x''(\omega)$ via

$$P_{\mathrm{abs}}(\omega) = \tfrac{1}{2}\omega\chi_x''(\omega)h_0^2. \tag{12}$$

Because of the overall isotropy, one can symmetrize over directions and use

$$\frac{P_{\mathrm{abs}}(\omega)}{h_0^2} = \frac{\omega^2 N}{2k_B T}\mathrm{Re}\int_0^\infty dt\, e^{i\omega t} A(t), \tag{13}$$

where $A(t) \equiv \langle \mathbf{m}(t) \cdot \mathbf{m}(0)\rangle/3$. Here, in the definition of $A(t)$, the subtraction terms were dropped as they do not contribute at finite frequencies.

III. NUMERICAL PROCEDURES

In this work, all computations were done for a 2D model of $300 \times 340 = 102\,000$ spins on a square lattice for $D_R/J = 1$. The free boundary conditions were used, although in the absence of long-range interactions for this rather large system the choice of boundary conditions is unimportant. The frequency-dependent absorbed power for this model at $T = 0$ was computed in Ref. [33] by pumping the system with ac fields of different frequencies. Here two numerical experiments were done.

In the first experiment the system was prepared at $T = 0$ by the energy minimization [11] starting from random directions of spins. It was subsequently pumped by an ac field with different large amplitudes h_0 at a fixed frequency $\omega/J = 0.075$

corresponding to the absorption maximum. The pumping routine was very long to allow nonlinearity, resonance saturation, and heating to develop. Equation (3) was used to compute the absorbed power.

The use of this method at elevated temperatures is problematic because fluctuations of $\mathbf{m}$ in a finite-size system easily dominate the response to a weak ac field $\mathbf{h}(t)$. Obtaining reliable results requires a prohibitively long computation to have these fluctuations averaged out, or it needs a large amplitude of the ac field h_0 that makes the response nonlinear.

Instead, we performed a second numerical experiment that used FDT to obtain the absorbed power $P_{\rm abs}(\omega, T)$ in the linear regime from the conservative dynamics of the system prepared by the Monte Carlo method at different temperatures. One computation of the conservative time evolution yields the results at all frequencies. As the averaging of fluctuations is a part of this procedure too, a long dynamical evolution is still needed.

Computing a long dynamical evolution of conservative systems requires an ordinary differential equation (ODE) solver that conserves energy. Many researchers prefer symplectic solvers that conserve energy exactly. However, the precision of popular symplectic solvers is low, with step error δt^3, while higher-order symplectic methods are cumbersome. In addition, these methods do not work for systems with single-site anisotropy. The mainstream general-purpose solvers such as Runge-Kutta 4 (with step error δt^5) and Butcher's Runge-Kutta 5 (with step error δt^6) do not conserve energy explicitly. Although the corresponding energy drift is small, it accumulates over large computation times.

This inconvenience can be cured by the recently proposed procedure of energy correction [37] that can be performed from time to time to bring the system's energy back to the expected target value. It consists of a (small) rotation of each spin towards or away from the respective effective field according to

$$\delta \mathbf{s}_i = \xi \mathbf{s}_i \times (\mathbf{s}_i \times \mathbf{H}_{{\rm eff},i}), \tag{14}$$

with ξ chosen such that the ensuing energy change $\delta E = \xi \sum_i (\mathbf{s}_i \times \mathbf{H}_{{\rm eff},i})$ has a required value. In the case of free conservative dynamics, one has $\delta E = E_0 - E$, where E_0 is the initial energy that must be conserved and E is the current energy that differs from E_0 because of the accumulation of numerical errors.

In our first experiment with the ac pumping, the situation with the energy conservation is the following. Computation of the absorbed energy $E_{\rm abs}(t) = \int_0^t dt' P_{\rm abs}(t')$ is robust. On the contrary, the energy of the system $E \equiv \mathcal{H}$ changes due to accumulation of numerical errors, even in the absence of pumping. If it changes significantly, then it begins to affect the absorbed power, and the whole computation breaks down. Thus, the energy of the system must be corrected such that the energy change $\Delta E(t) = E(t) - E_0$ equals the absorbed energy $E_{\rm abs}(t)$. This implies $\delta E = E_{\rm abs} - \Delta E$ in the energy-correcting transformation. In the case of pumping, this transformation should be done at moments of time when $\mathbf{h}(t) = 0$ to avoid the ac field making a contribution to the energy. In this work, it was done each time the period of the pumping was completed.

In our second experiment with the free conservative dynamics and FDT, the energy correction also was done with the fixed time intervals. In both experiments, the computation time was not limited by any accuracy factors.

The fifth-order Runge-Kutta method (the code can be found in the Appendix of Ref. [38]) with time step $\delta t = 0.1$ was used as the ODE solver. As the computational tool, we employed Wolfram *Mathematica* with compilation of heavy-duty routines into C++ code using a C compiler installed on the computer. Most of the computations were performed on our Dell Precision Workstation with 20 cores. The computation in the first numerical experiment was not parallelized. In the second experiment, to better deal with fluctuations, we ran identical parallel computations on the 16 cores available under our *Mathematica* license and averaged the results. We set J and all physical constants to 1 in all computations. In the presentation of the results the time t is given in the units of $\hbar/J$.

IV. NUMERICAL RESULTS

A. Pumping, nonlinearity, and heating

Pumping an isolated system for a long time increases its energy; that is, it produces heat. For realistically small ac amplitudes, $h_0/J \lll 1$, the computation times required to reach significant heating are prohibitively long, although the corresponding physical times may be quite short. Thus, to illustrate the effect of heating, we used extremely high amplitudes of the ac field. The results obtained at $h_0/J = 0.1$ are shown in Fig. 1. At long times, the energy change ΔE (which equals the absorbed energy $E_{\rm abs}$) reaches its maximal value corresponding to the total disordering of spins. In this state, resonances are saturated, and the absorbed power is close to zero. The dynamical spin temperature T_S defined by Eq. (4) becomes large. At $h_0/J = 0.03$ a very long computation allowed us to reach the spin temperature $T_S \approx 1.75J$.

The existence of the dynamic spin temperature T_S allows one to check whether the system is in thermal equilibrium during the energy-absorption process. In Fig. 2 the above results for $\Delta E(t)$ and $T_S(t)$ are parametrically replotted and compared with the dependence $\Delta E(T)$ computed by the Monte Carlo method. The parametric plots $\Delta E(T_S)$ for $h_0/J = 0.1$ and 0.03 perfectly coincide everywhere except for the smallest T_S, where the $h_0/J = 0.1$ curve bulges. For $h_0/J = 0.03$ the bulging also exists but is rather small. Both of these curves are in good accord with the Monte Carlo curve $\Delta E(T)$. This result strongly suggests that the absorbed energy goes into all modes and the equipartition of the energy is reached. It is remarkable that except for short times the system is in equilibrium even for the highest ac amplitude studied, $h_0/J = 0.1$. For realistically small ac amplitudes the equilibrium in the magnetic system should be complete.

The magnetic system being in equilibrium and having a particular spin temperature T_S during the absorption of the microwave energy allows one to set up a single differential equation for the temperature that describes the entire process. Expressing the time derivative of the energy in Eq. (3) via the heat capacity of the magnetic system C and the time derivative of the spin temperature, one obtains $\dot{\mathcal{H}}(t) = C\dot{T} = P_{\rm abs}(\omega, T)$,

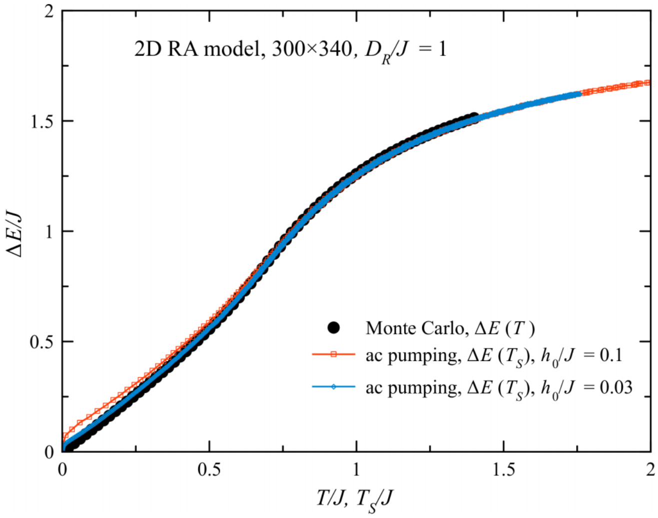

FIG. 2. Absorbed energy vs Monte Carlo temperature T and the spin temperature T_s.

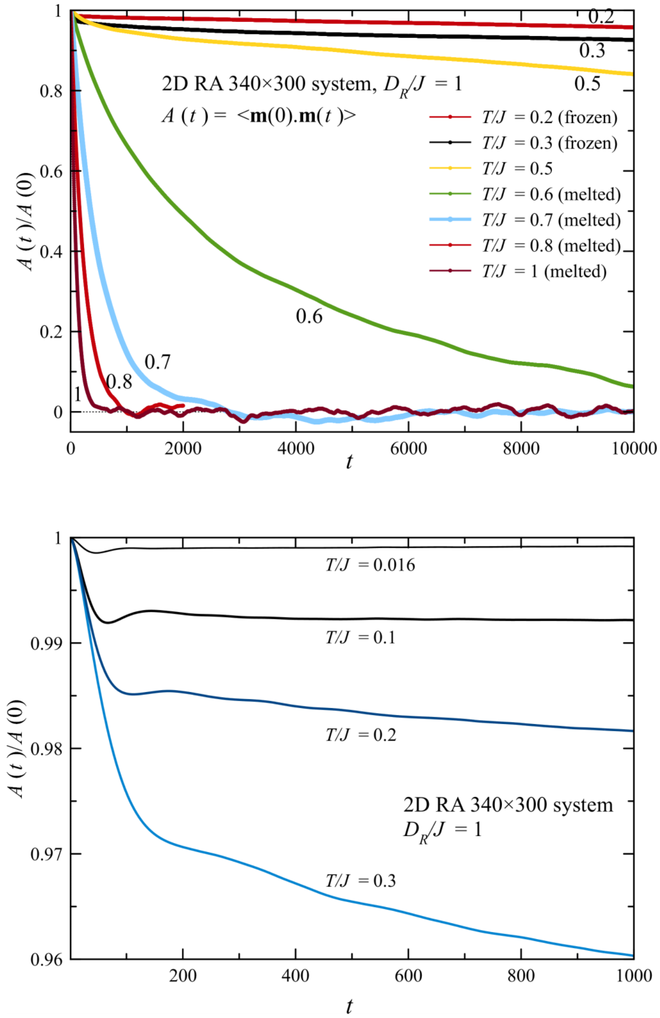

FIG. 3. Autocorrelation function of the total spin at different temperatures. Top: wide temperature range. Bottom: low temperatures.

which produces a simple equation for $T(t)$,

$$\dot{T} = \frac{P_{\text{abs}}(\omega, T)}{C(T)}. \tag{15}$$

Thus, it is sufficient to compute the heat capacity of the spin system $C(T)$ using the Monte Carlo method and compute the absorbed power $P_{\text{abs}}(\omega, T)$ at different temperatures to answer the question of heating and resonance saturation numerically. As realistic ac field amplitudes are rather small compared to the exchange even for a high-power microwave source, one can find $P_{\text{abs}}(\omega, T)$ in the linear regime with the help of the FDT. After the numerical solution of Eq. (15) is obtained, one can compute the absorbed energy by integration:

$$E_{\text{abs}}(t) = \int_0^t dt' P_{\text{abs}}[\omega, T(t')]. \tag{16}$$

Since the spin-lattice relaxation is rather fast, one can add the heat capacity of the lattice to that of the magnetic system. This must reduce heating and increase absorption. In addition the absorbed energy can flow from the magnetic particles to the dielectric matrix by heat conduction. However, this process is slow, and it cannot transfer a significant amount of energy during a short time of the microwave pulse.

B. Microwave power absorption by FDT

In the second numerical experiment, we ran conservative dynamical evolution of the states created using the Monte Carlo method at temperatures $T/J = 0.1 \div 1.0$ in steps of 0.1. Computation was performed in parallel using 16 processor cores until $t = t_{\max} = 50\,000$ in most cases. For the lowest temperature and for $T/J = 1$, the computation went up to $t = 100\,000$. From the computed dependence $\mathbf{m}(t)$ the autocorrelation function $A(t)$ entering Eq. (13) was obtained, and the results were averaged over the cores. Computations on our Dell Precision Workstation took 5 days for each temperature for $t_{\max} = 50\,000$. The normalized autocorrelation functions $A(t)/A(0)$ at different temperatures are shown in Fig. 3. These dependences have different forms at low and elevated temperatures that can be interpreted as a glassy transition or blocking due to barriers. Below $T/J \approx 0.5$, spins are frozen, and $A(t)$ is decreasing very slowly, apparently due to bunches of spins (IM domains) crossing anisotropy barriers via thermal agitation. Above this temperature, spins decorrelate with time fast, pointing to the meltdown of the magnetization pattern.

There is a fundamental unsolved question of whether the RA system can be described in terms of IM domains whose magnetic moments become blocked at low temperatures due to barriers or whether it is a true spin-glass state [39]. Since our focus is on the absorption of microwaves, we do not attempt to answer this question here. In Eq. (13) for the absorbed power, the low-frequency part of $A(\omega)$ is suppressed by the factor ω^2, so that the long-time physics is irrelevant for the absorption.

The main contribution to the absorbed power comes from the short-time part of $A(t)$ that is shown in the bottom panel of Fig. 3. There is an initial steep descent of $A(t)$ ending in a quasiplateau at low temperatures. That steep descent can be interpreted as being caused by dephasing of the precession of different IM domains in their potential wells. This preces-

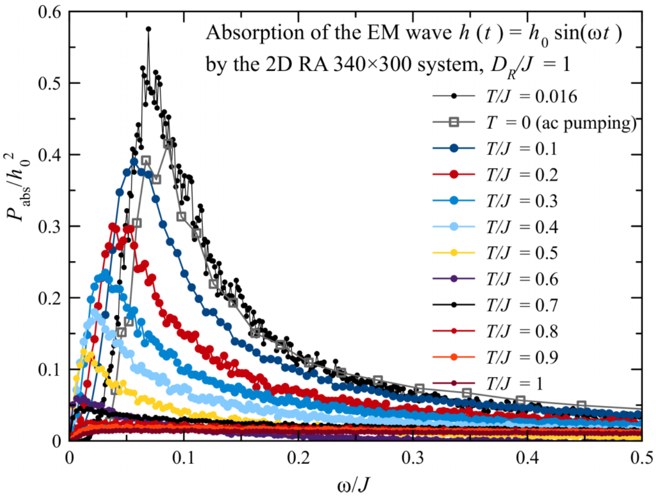

FIG. 4. Frequency dependence of the absorbed microwave power at different temperatures. The ac pumping result at $T = 0$ of Ref. [33] is shown for comparison.

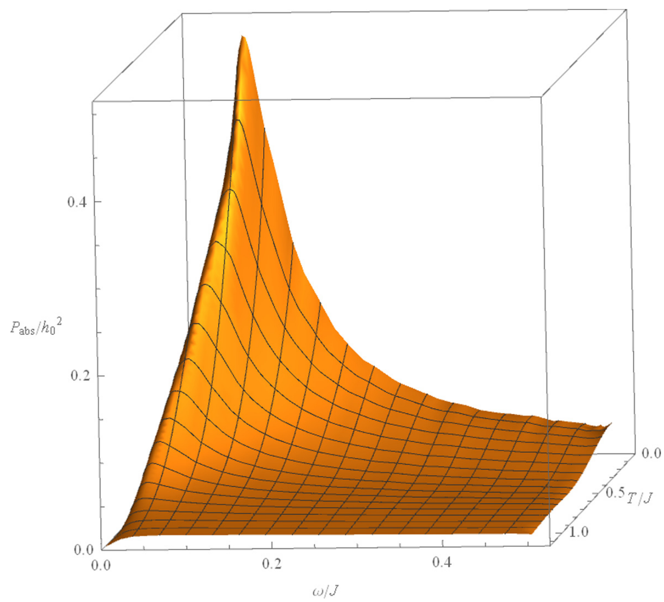

FIG. 5. Dependence of the absorbed microwave power on temperature and frequency.

sion with a quasicontinuous spectrum of frequencies is what ensures the absorption of the microwave power in a broad frequency range. As each IM domain remains precessing in its own potential well, $A(t)$ cannot change by a large amount. The latter requires flipping IM domains over the barriers, which happens at higher temperatures.

Making the Fourier transform of $A(t)$ and using Eq. (13), one obtains the absorbed power $P(\omega)$ that is shown in Fig. 4. One can see that the absorption curve gets depressed on increasing temperature, until the absorption peak vanishes when the glassy state melts. This is similar to what one observes in a system of independent resonating spins (or magnetic particles): The greater the thermal excitation of spins is, the less energy they absorb.

Absorption is maximal if the spins are in their ground states at $T = 0$. This case is also amenable to the FDT method if one uses a very small, but nonzero, temperature. Here the initial state was prepared using the Monte Carlo method at $T/J = 0.005$; then dynamical evolution was run. In the course of the evolution, the dynamical spin temperature T_S given by Eq. (4) changes just a bit. This change can be significant only at a very low temperature, and at room temperature it can be neglected. For instance, for the initial state with $T/J = 0.005$ we obtained $T_S/J = 0.0164$, but at $T/J = 0.5$ we got, for the average value of the spin temperature, $T_S/J = 0.504$. Thus, we used T_S instead of T in Eq. (13). One can see that the absorption at our lowest temperature (close to zero) is in fair agreement with the absorption data (also shown) obtained by the ac pumping at $T = 0$ [33]. However, the absorption due to the ac pumping near the maximum is somewhat depressed, apparently because of the partial saturation of resonances.

The method for computing the absorption based upon the FDT requires thermal equilibrium, that is, Boltzmann distribution over the energies of excitations. This is fulfilled at temperatures above freezing (blocking) of spins (see Fig. 3). At temperatures below freezing the FDT becomes applicable again as bunches of correlated spins are precessing at the bottoms of their energy minima and these low-energy excitations are in thermal equilibrium with each other. At such temperatures the processes of thermal activation over the barriers are slow and can be neglected in the frequency range relevant for the absorption. In the vicinity of freezing, the system is not at equilibrium, and the FDT becomes invalid. Here the spin-spin autocorrelation function $A(t)$ decays slowly, and it does not fit into the computation range $t < 10\,000$. Thus, the Fourier transforms of $A(t)$ that define the absorbed power are distorted for $T/J = 0.5$ and 0.6. Correspondingly, in Fig. 4 the curves for these temperatures are depressed at higher frequencies.

Figure 5 shows the same data presented in a 3D form as $P_{abs}(\omega, T)$ with the help of the approximation using Bézier functions implemented in Wolfram *Mathematica*. Bézier functions heal the depression of $P_{abs}(\omega, T)$ at $T/J = 0.5$ and 0.6, so it is not dramatic in this presentation.

Having built the approximation for $P_{abs}(\omega, T)$, one can plot the curves $P_{abs}(T)$ for different ω, as shown in Fig. 6. Since the maximum of $P_{abs}(\omega)$ shifts to the left on increasing T (see Fig. 4), the maximum in $P_{abs}(T)$ shifts to the left on increasing ω. The dependence $P_{abs}(T)$ was used in Eq. (15) for the spin temperature.

The shift of the absorption maximum to lower frequencies on increasing temperature that is seen in Figs. 4 and 5 can be explained in terms of the effective macroscopic anisotropy $K(T)$ that decreases with temperature. In conventional ferromagnets this effect is responsible for the narrowing of the hysteresis loop. For RA magnets the physics must be similar: Thermal fluctuations in the directions of spins tend to average out the anisotropy field acting on the spins. The frequency corresponding to the absorption maximum at $T = 0$ is defined by the microscopic RA averaged over IM domains [33]. On increasing temperature, the effective value of that anisotropy softens, which leads to the redshift in the absorption maximum. This effect has been observed in experiments with high-intensity microwaves in conventional ferromagnets, where it was attributed to spin-wave instabilities [40].

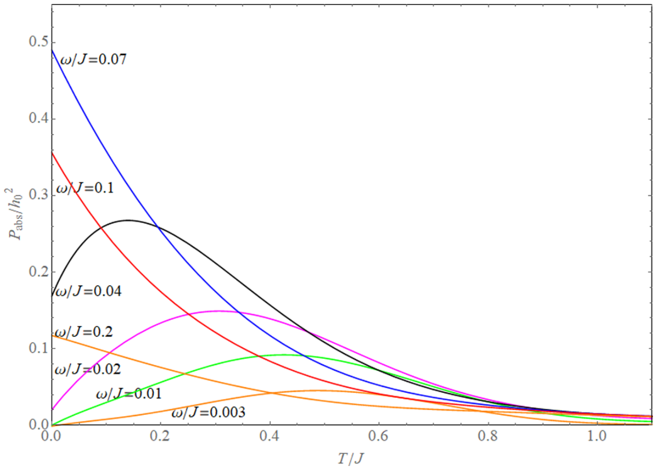

FIG. 6. Temperature dependence of the absorbed microwave power at different frequencies.

V. DISCUSSION

Previously, we showed [33] that RA magnets can be strong broadband absorbers of the microwave power. In this paper we have studied the temperature dependence of the power absorption. Our results answer two questions, both related to applications of microwave absorbers. The first one is a direct question of how the absorption by RA magnets depends on their temperature when they are heated by an independent source. It is answered by Fig. 4, which shows a decrease of the absorption on increasing temperature in a broad frequency range, making the system basically transparent for the microwaves at a sufficiently high temperature.

The second question is the response of the RA magnet to a microwave pulse of high power. We have shown that during the pulse, the spin system is in equilibrium and can be described in terms of the spin temperature. The typical times of spin-phonon transitions are much shorter than a microsecond, thus allowing the spin system to quickly equilibrate with the atomic lattice as well. On the contrary, the time required for the flow of heat out of the dielectric layer containing densely packed RA magnets is relatively slow. For a layer of thickness l, average mass density ρ, specific heat c, and thermal conductivity k it can be estimated as $t \sim (\rho c l^2)/k$. At l of the order of a few millimeters it is in the ballpark of a fraction of a second or longer, depending on the substrate. Thus, for microwave pulses that are shorter than that, the heat-conductivity mechanism is irrelevant.

A sufficiently strong pulse of microwave energy directed at such a layer from a close distance, having a duration in the range of microseconds to milliseconds, can greatly diminish the absorption capacity of the layer during the action of the pulse, making it transparent for a weak microwave signal coming from afar. During that time, if the layer is covering a metallic surface, microwaves in a broad frequency range would pass through it with the minimum absorption and would be reflected by the metal. If one wants to minimize this effect, the layer should be made of very densely packed metallic RA magnets electrically insulated from each other by a very thin dielectric coating. High thermal conductivity of such a system would greatly increase its cooling via heat conduction and would make its absorbing capabilities more resistant to high-power pulses of direct microwave energy.

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