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Model calculation for the susceptibility of a quantum spin glass

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We have studied the dynamic susceptibility in a quantum spin glass which is described by an Ising model in a transverse field, the latter introducing quantum tunneling in an otherwise classical problem. An effective single-spin Hamiltonian is obtained on the basis of an extant thermofield dynamic approach. The coupling to a dissipative heat bath is designed to affect thermal as well as quantal fluctuations and to yield Glauber kinetics in the absence of the transverse field. A resolvent expansion of the underlying time-development operator is then set up and the dynamic susceptibility calculated to leading order in perturbation theory. It is shown that the frequency of the peak of the susceptibility shifts towards higher values of ω, its amplitude is reduced and peaks are broadened as quantum effects strengthen. Our results are in qualitative agreement with the experimentally studied crystal-field-split system of $LiHo_xY_{1-x}F_4$ by Wu *et al.*

I. INTRODUCTION

The relaxational dynamics of quantum spin glasses is a problem of great contemporary interest.[1] The presence of disorder is already known to give rise to unusual time-dependent characteristics such as stretched-exponential decay of correlation functions and consequent non-Debye behavior of response functions.[2] In addition, quantum effects become important at very low temperatures, well below the glass transition, as different parts of the free-energy surface can be linked through tunneling. The additional quantum interactions are expected to cause further interesting dynamical effects. Recently, Wu *et al.*[3] have experimentally studied the simplest prototype of a quantum spin glass, viz., an Ising model (with disorder) plus a transverse coupling. The latter makes the system quantum mechanical and is physically realized by the application of an external magnetic field to a crystal-field-split system of $LiHo_xY_{1-x}F_4$ (see Ref. 3 for details). They have measured the imaginary component of the dynamic susceptibility $\chi''(\omega)$ as a function of both ω and the strength of the transverse field above and below the spin-glass transition temperature T_g, defined in the absence of the transverse field. The application of a transverse field radically affects the time scale of the Ising system's response. The frequency ω_p of the peak in the dynamic susceptibility $\chi''(\omega)$ increases by orders of magnitude as the strength of the transverse field is increased. The experimental data also suggest that the low-frequency tails of $\chi''(\omega)$ are greatly suppressed, indicating that quantum routes to relaxation affect the long-time dynamics of the system. Further, it is found that a longitudinal field primarily depresses the amplitude of the response and cannot account for the observed shifts in ω.

In this paper, we make a detailed comparison between the measured $\chi''(\omega)$ and the results obtained from a dynamical theory of the transverse Ising model, in the mean-field approximation. The relaxational dynamics is studied under the influence of a purely dissipative heat bath which is treated in the Markovian limit. The relaxation behavior of the model system in the absence of the transverse field is the same as that obtained from Glauber dynamics. The interplay of this with additional time dependence caused by quantum transitions due to the transverse field is analyzed in detail through the calculation of a spin-spin correlation function.

The paper is organized as follows. In Sec. II, the basic model Hamiltonian is discussed and its mean-field limit obtained. Various terms describing the interaction with the heat bath are also assessed. We then set up a resolvent expansion of the bath-averaged time-development operator which is relevant for the calculation of the correlation function. Next, the expressions for the correlation function and the generalized susceptibility are evaluated in Sec. III. Numerical results, their comparison with experiments, and concluding remarks are then presented in Sec. IV.

II. MATHEMATICAL FORMULATION

A. The Hamiltonian in mean-field approximation

The Hamiltonian describing a quantum spin glass composed of N interacting spins in a transverse field Ω may be written as[3]

$$H=\sum_{i,j=1}^{N} J_{ij}\sigma_i^z\sigma_j^z-\Omega\sum_{i=1}^{N}\sigma_i^x \ . \tag{1}$$

Here, spins i and j are connected by a random exchange J_{ij} and the σ's are Pauli spin matrices. The transverse field plays the role of an operator which mixes formerly pure eigenstates. The random interaction J_{ij} is assumed to be independently distributed according to a Gaussian

$$P(J_{ij})=\frac{1}{\sqrt{2\pi J^2}}\exp\left[-\frac{J_{ij}^2}{2J^2}\right] . \tag{2}$$

As a first step, we obtain an effective single-spin Hamiltonian within the framework of a mean-field theory (MFT) of quantum spin glasses. A systematic MFT for the model of a quantum glass defined by Eqs. (1) and (2) has been carried out in Ref. 4 using the thermofield dynamic approach and a short-time approximation for the dynamic self-interaction. here we follow that analysis. The effective single-spin Hamiltonian in the present case reads

$$H_S = -h\sigma^z - \Omega\sigma^x \ . \tag{3}$$

Here h is an effective field acting along the z-axis and is due to the nonzero spin-glass order parameter q,

$$h(\xi) = \tfrac{1}{2} J \xi \sqrt{q} \ , \tag{4}$$

where ξ is the excess static noise arising from the random interaction J_{ij}.[4] The mean-field equations for the local polarization $p(\xi)$ and the spin-glass order parameter q are [4]

$$p(\xi) = r(\xi)\tanh[\tfrac{1}{2}\beta h_0(\xi)] \tag{5}$$

and

$$q = \int_{-\infty}^{\infty} \frac{d\xi}{\sqrt{2\pi}} e^{-\xi^2/2}[p(\xi)]^2 \tag{6}$$

with

$$h_0(\xi) = \sqrt{\Omega^2 + h^2(\xi)} \tag{7}$$

and

$$r(\xi) = h(\xi)/h_0(\xi) \ . \tag{8}$$

The Hamiltonian of Eq. (3) describes the *reversible* dynamics of the system only. We now assume that the system described by H_S is placed in contact with a heat bath. In the absence of the tunneling term, i.e., $\Omega=0$, it is customary to imagine that the dynamics arises from additional coupling terms to the heat bath which are off diagonal in the representation in which σ^z is diagonal. This is then in the spirit of the kinetic Ising model of the Glauber type. The physical meaning of the Glauber terms is evident. They cause spontaneous spinflips, which, in the context of our system, mimic thermally activated jumps between the ground-state doublet. In the case of which $\Omega \neq 0$, we could still imagine the heat-bath interactions to be of the Glauber type. But that would be tantamount to taking into account only the classically activated jump processes. In reality, we expect the heat bath to induce not only thermal fluctuations of the above kind but quantum fluctuations as well, leading to incoherence in tunneling which is otherwise a coherent phenomenon. It is important here to remember that we would like to recover the Glauber mechanism in the limit $\Omega=0$. The point is that, when $\Omega \neq 0$, the appropriate quantization axis is neither z nor x, but somewhere in between. The simplest coupling to the heat bath should then involve an operator that is strictly off diagonal in the new representation of the quantization axis and would lead to correct limits when $\Omega=0$. Hence it is necessary to perform a rotation in the "spin space" of the subsystem by an angle $\theta=\arctan(\Omega/h)$ in the x-z plane around the y axis (see further remarks below). Motivated by our preceding comments, we generalize the Hamiltonian as in (3) to

$$H_0 = H_S + H_I + H_B \ , \tag{9}$$

where H_I describes the interaction between the spin subsystem and the heat bath. In accordance with our stated objective, we assume the following type of interaction:

$$gb\left[\frac{h}{h_0}\sigma^x + \frac{\Omega}{h_0}\sigma^z\right] . \tag{10}$$

In (10), b is an operator which acts on the Hilbert space of the heat-bath Hamiltonian H_B and g is a multiplicative coupling constant. The specific form of the interaction is chosen so as the guarantee that, in the rotated frame in which H_S is diagonal, the coupling with the heat bath is purely off diagonal [Eq. (16d) below]. The exact nature of the operator b will not be specified here; suffice it to say, however, that the coupling term is expected to yield Glauber kinetics for the underlying Ising model if the tunneling term were absent.[5] On the other hand, if Ω is nonzero, the term proportional to it will lead to incoherent effects on tunneling. The fact that the ratio of the coupling terms in (10) (proportional to σ^x and σ^z respectively) is taken as h/Ω does not seem to cause any loss of generality.

We remark here that a stochastic formalism[6] can also be adopted to study the quantum effects of tunneling. Work along these lines has been done by us earlier in the context of the proton-glass problem,[7(a)] where the subsystem is treated quantum mechanically, while the surrounding heat bath is handled as a classical stochastic reservoir. However, we find that a calculation of the dynamic susceptibility in this formalism does not exhibit the characteristic features observed by Wu *et al.*[3]

B. Averaged time-development operator

We first note that the expression for susceptibility due to an oscillatory magnetic field applied along the z axis is[6]

$$\chi(\omega) = \tfrac{1}{2}\beta \lim_{\substack{\delta\to 0 \\ s\to -i\omega+\delta}} \left[\frac{1}{s} - 4\tilde{c}(s)\right] , \tag{11}$$

where $\tilde{c}(s)$ is the Laplace transform of the correlation function $c(t)$ defined as

$$c(t) = \langle \sigma_z(0)\sigma_z(t)\rangle \ . \tag{12}$$

Here the angular brackets denote the appropriate quantum and statistical average. The quantity s is related to the applied frequency ω, $s=-\omega i+\delta$, δ being a small real-valued parameter, and β is the inverse temperature. Explicitly, $c(t)$ can be expressed as

$$c(t) = \frac{1}{Z_0}\mathrm{Tr}[e^{-\beta H_0}\sigma^z(0)e^{iH_0 t}\sigma^z(0)e^{-iH_0 t}] \ , \tag{13}$$

where H_0 is the total Hamiltonian as in Eq. (8) and Z_0 is the corresponding partition function.

As discussed in Sec. II A, we diagonalize the subsystem Hamiltonian H_S by performing a rotation U^R in the spin space of the subsystem by an angle $\theta=\arctan(\Omega/h)$ around the y axis:

$$U_Y^R = e^{-i(\theta/2)\sigma^y} \, . \tag{14}$$

In the rotated frame, the correlation function reads

$$c(t) = \frac{1}{Z_0} \mathrm{Tr} \left[e^{-\beta \tilde{H}_0} \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] \times e^{i\tilde{H}_0 t} \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] e^{-i\tilde{H}_0 t} \right] , \tag{15}$$

where

$$h_0 = \sqrt{h^2 + \Omega^2} \, , \tag{16a}$$

$$\tilde{H}_0 = \tilde{H}_S + \tilde{H}_I + H_B \, , \tag{16b}$$

$$\tilde{H}_S = h_0 \sigma^z \, , \tag{16c}$$

$$\tilde{H}_I = g b \sigma^x \, , \tag{16d}$$

g and b having been defined as in Eq. (10). Assuming that the subsystem is coupled weakly to the heat bath, we can factorize the density matrix and write the correlation function as

$$c(t) = \frac{1}{Z_S} \mathrm{Tr}_S \left\{ e^{-\beta \tilde{H}_S} \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] \mathrm{Tr}_B \left\{ \exp \frac{-\beta H_B}{Z_B} \left[U(t) \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] \right] \right\} \right\} , \tag{17}$$

where $U(t)$ is the time-development operator. The Laplace transform of $c(t)$ reads

$$\tilde{c}(s) = \frac{1}{Z_S} \mathrm{Tr}_S \left\{ e^{-\beta H_S} \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] \times \left\{ [\tilde{U}(s)]_{\mathrm{av}} \left[\frac{h}{h_0}\sigma^z + \frac{\Omega}{h_0}\sigma^x \right] \right\} \right\} , \tag{18}$$

where $[\tilde{U}(s)]_{\mathrm{av}}$ denotes the Laplace transform of the time-development operator averaged over the bath degrees of freedom. The effects of dynamics are contained in this bath-averaged operator.

As discussed extensively in Ref. 6, it is the physics of a given problem that decides the nature of the time-development operator. In the present context, we have adopted the system-plus-reservoir approach in order to give a proper treatment of the incoherent tunneling term, and systematically "project out" the bath degrees of freedom. This can be most conveniently achieved by writing a resolvent expansion of $[\tilde{U}(s)]_{\mathrm{av}}$ in which the interaction term H_I is treated perturbatively. As discussed in detail in Ref. 6, such an expansion yields the following general expression for $[\tilde{U}(s)]_{\mathrm{av}}$:

$$[\tilde{U}(s)]_{\mathrm{av}} = [s - iL_S + \tilde{\Sigma}(s)]^{-1} \, , \tag{19}$$

where L_S is the Liouville operator associated with the spin Hamiltonian H_S in Eq. (3) and $\tilde{\Sigma}(s)$ is the so-called relaxation matrix, to be specified below. While it is possible to evaluate $\tilde{\Sigma}(s)$ to arbitrary orders in perturbation theory, it suffices for the purpose of obtaining Markovian dynamics to use the expansion to second order in H_I, which yields

$$\tilde{\Sigma}(s) = \left\{ L_I \frac{1}{s - iL_S - iL_B} L_I \right\}_{\mathrm{av}} \, . \tag{20}$$

In the next section, we present the details of the calculation of the correlation function and the generalized susceptibility.

III. CORRELATION FUNCTION AND GENERALIZED SUSCEPTIBILITY

Having set up the formalism for evaluating the bath-averaged time-development operator, we now proceed to calculate the frequency-dependent susceptibility. First, we evaluate the different components of $[\tilde{U}(s)]_{\mathrm{av}}$. The matrix elements of a Liouville operator L_S can be expressed in terms of the matrix elements of the corresponding subsystem Hamiltonian as[6]

$$(\nu\mu | L_S | \nu'\mu') = [\delta_{\mu\mu'} \langle \nu | H_S | \nu' \rangle - \delta_{\nu\nu'} \langle \mu' | H_S | \mu \rangle] \, . \tag{21}$$

Here we denote the "matrix elements" of L by parentheses. These are labeled by four indices, just as the elements of H_S are labeled by two. It is now clear from Eq. (21) that the evaluation of the matrix elements of $[\tilde{U}(s)]_{\mathrm{av}}$ involves the inversion of a matrix which in the present problem has a dimension of 4×4 as the spin operator σ^z is only a two-valued operator. The first component of this matrix, $(\mu\nu | L_S | \mu\nu)$, in a rotated frame of reference, has the following 4×4 representation in which H_S is diagonal:

$$\begin{bmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & -h_0 & 0 \\ 0 & 0 & 0 & h_0 \end{bmatrix} , \tag{22}$$

where the rows and columns labeled by $|\mu\nu)$ take the values $|++)$, $|--)$, $|+-)$, and $|-+)$, respectively, and $h_0 = \sqrt{h^2 + \Omega^2}$.

The next step is the evaluation of the relaxation matrix, for which we proceed along lines similar to our earlier work in the context of the proton-glass problem.[7(b)] Treating the heat bath in the Markovian approximation, we had earlier found that

$$\tilde{\Sigma}(s) \approx \tilde{\Sigma}(0) = \int_0^\infty dt \{ L_I \exp[(L_S + L_B)t] L_I \} \ . \quad (23)$$

We further assume that the influence of the heat bath is purely dissipative; i.e., the relaxation matrix $\tilde{\Sigma}(s=0)$ is real. As observed in Ref. 7(b), all the elements of $\tilde{\Sigma}(s=0)$ can be expressed in terms of certain bath correlation functions. These correlation functions are not evaluated explicitly, but simply parametrized in terms of a phenomenological relaxation rate λ [see Ref. 7(b) for details].

In order to obtain the matrix of $[\tilde{U}(s)]_{av}$, we have to invert the matrix M:

$$\begin{bmatrix} s+\frac{\lambda}{2}(1-p_0) & -\frac{\lambda}{2}(1-p_0) & 0 & 0 \\ -\frac{\lambda}{2}(1+p_0) & s+\frac{\lambda}{2}(1+p_0) & 0 & 0 \\ 0 & 0 & s+\lambda+ih_0 & -\lambda \\ 0 & 0 & -\lambda & s+\lambda-ih_0 \end{bmatrix}, \quad (24)$$

where $p_0 = p_+^0 - p_-^0$ is the net polarization in the rotated frame and $p_\pm^0$ are the respective probabilities for the states $|\pm\rangle$. The matrix M is easily invertible and yields the following bath-averaged time-development operator $[\tilde{U}(s)]_{av}$ for the system we consider:

$$[\tilde{U}(s)] = \begin{bmatrix} \frac{1}{s(s+\lambda)} \begin{bmatrix} s+\lambda p_+^0 & \lambda p_-^0 \\ \lambda p_+^0 & s+\lambda p_-^0 \end{bmatrix} & \begin{matrix} 0 & 0 \\ 0 & 0 \end{matrix} \\ \begin{matrix} 0 & 0 \\ 0 & 0 \end{matrix} & \frac{1}{s(s+2\lambda)+h_0^2} \begin{bmatrix} s+\lambda-ih_0 & -\lambda \\ -\lambda & s+\lambda+ih_0 \end{bmatrix} \end{bmatrix} . \quad (25)$$

We may parenthetically comment here that the matrix of $[U(s)]_{av}$ in (25) is only slightly different from the one obtained from a purely stochastic formalism.[7(a)] This small difference is, however, crucial for properly accounting for quantum effects, characterized, for instance, by the off-diagonal nature of the underlying density operator of the system.

Using these matrix elements the relevant correlation function can be obtained from the expression (18). The susceptibility is then calculated by using the correlation function in Eq. (11) and averaging over the distribution of local polarization p. In view of Eq. (5), this is equivalent to averaging over the excess static noise field ξ which appears in the effective single-spin Hamiltonian (3). Since the experimental results of Wu *et al.* involve a study of the imaginary component of the ac susceptibility $\chi''(\omega)$, we calculate the same to facilitate a comparison with our theory. We finally express $\chi''(\omega)$ as

$$\chi''(\omega) = \frac{1}{4\pi} \int_{-1}^{1} dp \ W(p) \chi''(\omega, p) \ , \quad (26)$$

where the susceptibility for a given configuration of local polarization p is given by

$$\chi''(\omega,p) = \frac{\beta}{4} \left[\frac{\omega\lambda}{\omega^2+\lambda^2} \left[\frac{h^2}{h_0^2} - p^2 \right] + \frac{2\omega\lambda h_0^2}{(h_0^2-\omega^2)^2+4\omega^2\lambda^2} \frac{\Omega^2}{h_0^2} \left[1 - \omega \frac{p}{h} \right] \right] . \quad (27)$$

The averaged distribution for the local polarization can be defined as[8]

$$W(p) = \frac{1}{N} \sum_i [\delta(p - \langle \sigma_i^z \rangle)]_{av} = [(p - \langle \sigma_{iz} \rangle)]_{av} \ , \quad (28)$$

where $\langle \cdots \rangle$ denotes a thermal average and $[\cdots]_{av}$ a combined average over the distribution of random interactions. The Edwards-Anderson order parameter q is thus given by the second moment of $W(p)$. In terms of $p(z)$, the local polarization function [Eq. (28)] becomes

$$W(p) = \int D\xi \, \delta(p - p(\xi)) \ . \quad (29)$$

Using the well-known relation $\delta[f(\xi)] = \delta(\xi - \xi_0)/f(\xi_0)$, ξ_0 being the solution of $f(\xi)=0$, we obtain for the case $\sigma^z = \pm\frac{1}{2}$

$$W(p) = \frac{4}{\beta J \sqrt{2\pi q}} e^{-\xi_0^2/2} \times \left[\frac{h^2(\xi)}{h_0^2(\xi)} + \frac{2\Omega^2}{\beta h(\xi) h_0^2(\xi)} p - p^2 \right]^{-1} , \quad (30)$$

where $\xi_0 = \xi_0(p)$ is the inverse function of $p(\xi)$ which satisfies Eq. (5) for p in the interval $[-1, +1]$.

IV. RESULTS, DISCUSSION, AND CONCLUSION

The result of a numerical evaluation of the probability distribution of the local polarization is presented in Fig. 1 where $W(p)$ is plotted against p for a fixed value of J ($=10$) and various values of Ω ($=5.0$, 3.5, and 1.0). It can be seen from Fig. 1 that, if one varies the transverse field Ω, the shape of $W(p)$ changes from a peak structure at $p=0$ to that at $p=1$ as the value of Ω is decreased from 5.0 to 1.0. The variation in Ω is achieved experimentally by a magnetic field applied perpendicular to the c axis of the $LiHo_xY_{1-x}F_4$ system. It is also observed from Fig. 1 that, for certain values of the ratio α of Ω and J (~ 0.35), $W(p)$ is more or less a constant over a wide range of p values except near the end points $p = \pm 1$. From Eq. (30) and Fig. 1 it is clear that $W(p)$ can be ap-

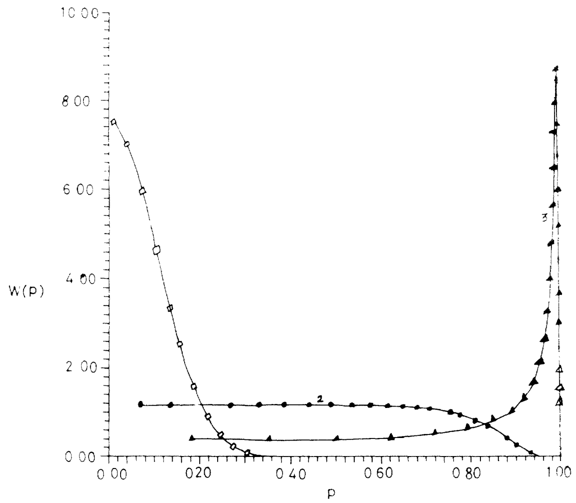

FIG. 1. Variation of the local polarization function $W(p)$ [given by Eq. (30) in the text] as a function of the polarization p for a fixed disorder $J=10$ and three different values of the transverse field Ω: curve 1, $\Omega=5.0$ (squares), curve 2, $\Omega=3.5$ (circles), and curve 3, $\Omega=1.0$ (triangles).

proximated by a constant function for values of α close to 0.35 over a wide range of p. This scheme enables us to compute $\chi''(\omega)$ easily, which facilitates further analysis.

Replacing $W(p)$ in Eq. (26) by its value at $p=0$, we obtain after carrying out the remaining integration the variation of $4\chi''(\omega)/\beta$ with respect to ω in the absence of quantum effects with a relaxation rate λ $(=0.9)$, as depicted in Fig. 2. As mentioned earlier, when quantum effects are absent, our model reduces to the Glauber model for Ising kinetics in the mean-field limit. Also shown in the figure are curves corresponding to different values of the ratio α with the same relaxation rate λ; typical values of α which have been selected are 0.3, 0.38, and 0.45. These values of α imply that the quantum effects are moderate in strength as compared to the disorder effects. It should also be noted that, for these values of α, our approximation of replacing $W(p)$ by a constant is a reasonable one. Examining the dynamic susceptibility response around its peak value (Fig. 3) clearly indicates a shift in the peak frequency ω_p as α is increased. There is also a reduction in the amplitude of $\chi''(\omega)$ and a broadening of peaks as quantum effects are enhanced. These features are qualitatively similar to the experimental observation made by Wu *et al.* in the $LiHo_xY_{1-x}F_4$ ferromagnet at temperatures below the spin-glass transition temperature T_g defined in the absence of the transverse field. We have checked that these features are not reproducible when λ is very small, i.e., heat-bath-induced effects are negligible. Hence we conclude that heat-bath coupling of the kind assumed by us is necessary to bring out the essential features of the quantum and relaxational dynamics of the system studied by Wu *et al.*[3]

Summarizing, we have considered in this paper the dynamics of a quantum spin glass described by the Ising model plus a transverse field and Gaussian disorder, to analyze the recently studied ferromagnetic system $LiHo_xY_{1-x}F_4$ (Ref. 3). Quantum effects similar to tunneling are introduced by the transverse field Ω which causes a mixing between the ground and the excited states of the crystal-field-split system. The relaxation dynamics has been studied by coupling the quantum subsystem to a purely dissipative heat bath. The coupling has been chosen to induce thermal as well as quantum fluctuations and the Glauber mechanism is recovered in the

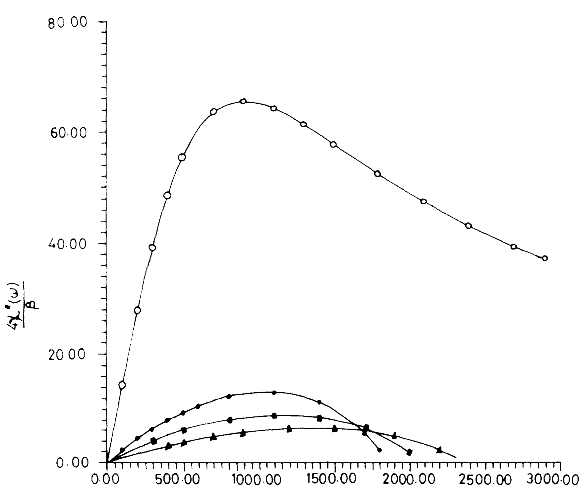

FIG. 2. Variation of the imaginary component of the dynamic susceptibility $\chi''(\omega)$ as a function of ω, the frequency of the applied ac field, in the absence of quantum effects ($\Omega=0$, open circles) and three different values of α $(=\Omega/J)$, viz., 0.3 (circles), 0.38 (squares), and 0.45 (triangles). $W(p)$ can be approximated by a constant for these values of α. $\chi''(\omega)$ exhibits a peak at $\omega=900$, 1000, 1200, and 1400 for $\alpha=0.0$, 0.3, 0.38, and 0.45, respectively.

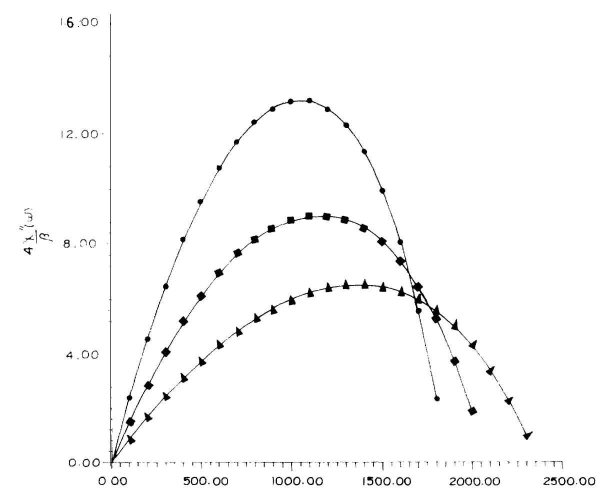

FIG. 3. Enlarged view of the variation of $\chi''(\omega)$ (as in Fig. 2) around its peak value. A shift of the peak of $\chi''(\omega)$ to the right, a decrease in its amplitude, and a broadening of the peak with increase in the strength of quantum effects is clearly observed for the chosen values of α. These observations are in agreement with the experimental measurement of $\chi''(\omega)$ on the crystal-field-split system of $LiHo_xY_{1-x}F_4$ by Wu *et al.* (Ref. 3).

absence of the tunneling term. An effective single-spin Hamiltonian has been obtained within a thermofield dynamic approach, which is applicable to the spin-glass phase above the instability surface. We have then computed the dynamic susceptibility $\chi''(\omega)$ as a function of both ω and α, α being the ratio of the transverse field Ω to the disorder J. Our main conclusion is that for a certain range of values of the ratio $\alpha(\sim 0.3-0.45)$ the frequency of the peak of the dynamic susceptibility shifts towards higher values of ω as α is increased. The strength of the susceptibility, on the other hand, is reduced and the peak is broadened as quantum effects are strengthened. The experimental results of Wu *et al.*[3] are in qualitative agreement with our calculations for the selected range of α. Most recently, Wu *et al.*[9] have performed nonlinear susceptibility measurements at $T=0$ and have observed a first-order phase transition driven by the transverse field. Some of these results can be explained within our theoretical framework and will be presented separately.

ACKNOWLEDGMENTS

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