SPIN-GLASS DYNAMICS IN THE TWO-DIMENSIONAL ISING SYSTEM $Rb_2Cu_{1-x}Co_xF_4$

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In the dynamics of random magnetic systems, the dimensionality of the magnetic lattice plays a crucial role. Two-dimensional ($d=2$) Ising spin glasses do not order until $T=0$ K, and accordingly the dynamics at long length scales is anticipated to be dominated by thermal activation over energy barriers.[1-3] Upon approaching the transition, this activated dynamical behavior involves a much faster divergence of relaxation times than ordinary critical slowing down at a conventional phase transition. The present study provides an experimental verification of these predictions in an actual $d=2$ system.

The experiments are performed on a single crystal of $Rb_2Cu_{1-x}Co_xF_4$, with $x=0.218$, which provides an excellent example of a random-bond $d=2$ spin glass with competing Ising-like ferromagnetic and antiferromagnetic interactions.[4] Results for the real and imaginary parts of the magnetic susceptibility in the frequency range 0.3-50 kHz are shown versus the temperature in Fig. 1. The data have been obtained by use of conventional mutual-inductance techniques. Noteworthy features are the strong frequency dependence of the maximum of $\chi'(T)$, the extremely weak frequency variation of $\chi''(T)$ above the freezing temperature, and the crossing near 4 K of the $\chi''(T)$ curves at various frequencies.

To analyze the data, we assume a specific time dependence for the decay of the spin-autocorrelation function $q(t)=\langle S_z(0)S_z(t)\rangle$, and subsequently calculate the frequency dependence of the susceptibility by numerical integration from the relation

$$\chi(\omega) = -\chi_0\int_0^\infty \left(e^{-i\omega t}\frac{d}{dt}q(t)\right)dt \quad , \tag{1}$$

in which χ_0 is the isothermal susceptibility. First, we consider the stretched-exponential decay

$$q(t) = \exp\left[-(t/\tau_c)^\beta\right] \quad , \tag{2}$$

with β a temperature-dependent exponent, and τ_c a characteristic relaxation time. Fits of Eqs.(1) and (2) to the data, with χ_0, β, and τ_c as adjustable parameters, turn out to be of good quality in the temperature range from 7 down to 3.4 K. The value of the exponent β, presented in Fig. 2 versus the temperature, is found to increase gradually from $\beta=0.06$ at 7 K to a

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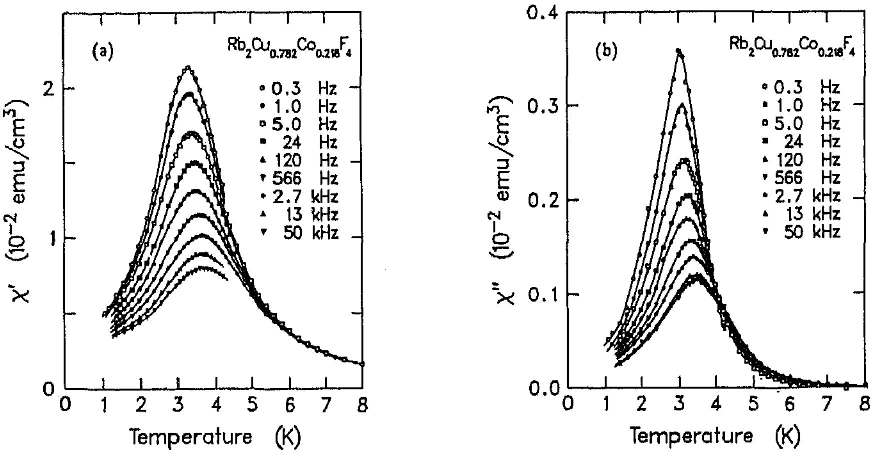

Fig.1. (a) Real part $\chi'(\omega,T)$ of the linear susceptibility vs the temperature for $Rb_2Cu_{0.782}Co_{0.218}F_4$; (b) Same as (a), but imaginary part $\chi''(\omega,T)$.

constant value $\beta = 0.096$ below 4 K. These very small values evidence an extremely slow decay of $q(t)$, much slower than in the case of short-range $d = 3$ spin glasses, for which numerical simulations[5] and experiments[6] both give $\beta \approx 0.3$ just above the spin-glass transition. As to the time parameter in Eq.(2), τ_c is found to increase from 10^{-9} s at 7 K to 10^{+3} s at 3.4 K, an extremely fast increase covering many more decades than the experimental time window of $3\times10^{-6} < \omega^{-1} < 0.5$ s. Accordingly, measurements at low and high temperatures probe, respectively, the short-time ($\omega^{-1} < \tau_c$) and long-time ($\omega^{-1} > \tau_c$) behavior of $q(t)$. The cross-over occurs near 4 K, where τ_c^{-1} falls in the middle of the experimental frequency window. As it appears from Fig. (2), the *short*-time decay of $q(t)$ can be quite well represented by Eq.(2), with a temperature-independent exponent $\beta = 0.096$.

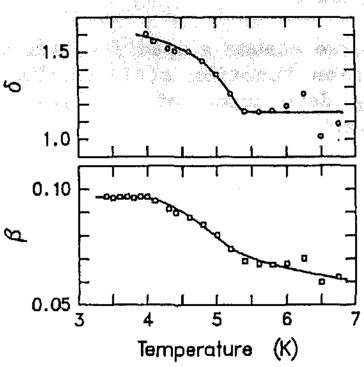

Fig.2. Exponents β and δ, occurring in Eqs.(2) and (3), vs the temperature.

The smallness of β as well as its anomalous temperature dependence above 4 K suggest a slower decay at long times than given by Eq.(2). Therefore, we have also considered the exponential-logarithmic time decay

$$q(t) = \exp\left[-\left(\frac{\ln(t/\tau_0)}{\ln(\tau_c/\tau_0)}\right)^{\delta}\right] , \tag{3}$$

with τ_0 a microscopic relaxation time, for which we take the single-spin flip time 10^{-13} s. Equation (3) is an empirical relation found to describe[7] the long-time tail of the decay in $d=3$ random-field Ising systems, for which the dynamics is expected to be closely similar to $d=2$ Ising spin glasses. Fits to the data of Eq.(1) with Eq.(3) substituted indeed turn out to be better than those of the stretched-exponential decay above 4.6 K ($\chi^2 \approx 0.3$ compared to $\chi^2 \approx 0.6$), but worse below this temperature ($\chi^2 \approx 2$ compared to $\chi^2 \approx 1$). In the range 5–7 K, the exponent δ, presented in Fig. 2 versus the temperature, appears to be constant ($\delta = 1.15 \pm 0.05$), implying an almost algebraic ($\delta = 1$) decay of $q(t)$. Thus, the *long*-time tail of $q(t)$ is well described by Eq.(3) with a constant δ. At lower temperatures, the fits gradually deteriorate, and the value of δ steadily increases. It is of interest to compare the present results with those for $d=3$ Ising spin glasses, where the decay $q(t)$ is found[6,7] to be adequately described by a combination of algebraic and stretched-exponential decays, however in a reversed order compared to the $d=2$ case and with different temperature-dependent exponents.

As to the temperature dependence of the characteristic relaxation time τ_c, results for $\ln(\tau_c/\tau_0)$ are plotted double-logarithmically as a function of the temperature in Fig. 3. Relaxation times derived from the stretched-exponential fit have also been included. A very rapid increase is observed with decreasing temperature. In the temperature range where δ and β are

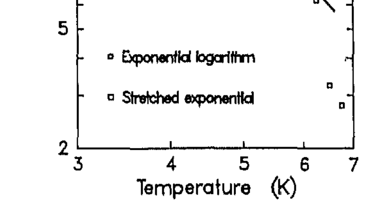

Fig.3. Double-logarithmic plot of $\ln(\tau_c/\tau_0)$ vs the temperature, with τ_c derived from Eq.(2) (squares) and Eq.(3) (circles).

constants, straight lines have been drawn through the respective data sets, corresponding to the dependence

$$\ln(\tau_c/\tau_0) \propto T^{-1-\psi\nu} \quad . \tag{4}$$

This behavior has been predicted[8] to occur in $d=2$ spin glasses on the grounds that the dynamics is governed by thermally activated relaxation over energy barriers whose height B scales as $B \propto \xi^{\psi} = T^{-\psi\nu}$, with ξ the spin-glass correlation length, and ψ and ν critical exponents. From the slope of the solid lines in Fig. 3, we find $\psi\nu = 1.9 \pm 0.2$ for the stretched-exponential decay, and $\psi\nu = 1.6 \pm 0.2$ for the exponential- logarithmic case. An analysis in terms of the Cole-Cole formalism yields $\psi\nu = 2.2 \pm 0.2$,[9] while a dynamic-scaling analysis of all $\chi''(\omega,T)$ data results in $\psi\nu = 2.2 \pm 0.3$.[9] Thus, irrespective of the particular analysis, compelling evidence for activated dynamics according to Eq.(4) is found. Averaging, somewhat arbitrarily, over the $\psi\nu$ values from the various methods, we find $\psi\nu = 2.0 \pm 0.3$. Taking $\nu = 2.4 \pm 0.3$, as derived from static susceptibility measurements, we then arrive at $\psi = 0.8 \pm 0.2$, which is within the theoretical limits $0 \leq \psi \leq d-1$,[8] and, within errors, equals $\psi = 1$ from a recent numerical calculation.[10]

Finally, we calculate the distribution of relaxation times $g(\tau)$ from the inverse Laplace transform of $q(t)$. For very small values of β, i.e., $\beta|\ln(t/\tau_c)| \ll 1$, $g(\tau)$ can be written as a series expansion in β,[11]

$$g(\tau) = (\beta/e)\exp[-\tfrac{1}{2}\beta^2\ln^2(\tau/\tau_c)][1+O(\beta^2)] \quad , \tag{5}$$

with τ_c the median relaxation time. To leading order, Eq.(5) represents a Gaussian distribution on a $\ln\tau$ scale. Because of the smallness of β, the distribution is extremely broad, spanning over 9 decades in time (full width at half maximum), even far above the freezing temperature.

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