Internal Field Distributions of Mesoscopic Spin Glasses

P. G. N. de Vegvar and T. A. Fulton
AT&T Bell Laboratories, Murray Hill, New Jersey 07974
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Magnetoconductance fluctuations are used to extract information about frozen spin configurations in submicrometer sized *Cu*Mn wires with Mn concentrations of 320 and 1000 atomic ppm (parts per 10^6). Measurements with the magnetic field parallel to the wires separate the spin and orbital field scales. At low temperatures the configurations distort reversibly over a field scale on the order of the bulk exchange field H_g. The data are used to map out a sample specific internal field distribution to over $30H_g$.

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Over distances comparable to the electron dephasing length L_ϕ the electrical conductance is sensitive to the specific scattering potential imbedded in the particular sample being measured [1]. Relying on this fact, quantum transport at low temperatures may be utilized as a probe of spin configurations in disordered magnetic systems [2,3]. When the specimen size L approaches L_ϕ, the low frequency four-point resistance R is a combination of interlead transmission coefficients and contains information about how the frozen spins elastically scatter conduction electrons. Our earlier work demonstrated how the time reversal symmetry-breaking property of frozen magnetic systems like spin glasses may be exploited to extract a magnetofingerprint of each particular static spin configuration directly from a linear combination of conductance fluctuations [4]. We previously studied *Cu*Mn spin glass wires and rings by measuring such "spin fingerprints" only up to fields on the order of the bulk exchange field $H_g = k_B T_g/\mu$ (T_g the bulk zero field spin freezing temperature, μ the Mn ion magnetic moment). To learn how static spin configurations transform over the H-T plane, we performed spin fingerprint measurements in several geometries and up to fields as high as $34H_g$ on *Cu*Mn wires of two concentrations. From experiments on wires oriented parallel to the applied field we find a field scale for reversible spin distortion of order H_g. When one pictures a spin glass in terms of a distribution of molecular fields, then the parallel field data can be used to determine this distribution. To our knowledge, this is the first measurement of such a function itself as opposed to merely its characteristic width [5,6].

The fabrication of the *Cu*Mn films and wires was discussed in previous work [4] and involved UHV Ar ion sputtering, e-beam lithography, and liftoff techniques. The linewidths W were in the range 800–850 Å with film thicknesses 300–400 Å, barely entering the regime where finite-size effects suppress T_g [7]. All the films displayed the same characteristic zero field $R(T)$, namely, Kondo-like for $T > T_g$ and nearly temperature independent for $T < T_g$, with an elastic mean free path at 4.2 K of 110–230 Å [4]. Bulk values for T_g and H_g are given in Table I.

The spin fingerprint of a frozen spin system where $\mathbf{S}_i(-H) \neq -\mathbf{S}_i(H)$ utilizes the resulting broken Onsager-Büttiker (OB) symmetry [8] for a four-point resistance: $R_{IV}(H) \neq R_{VI}(-H)$. Here I (V) denotes a pair of current (voltage) measuring leads. Therefore the OB antisymmetric and field symmetric combination of magnetoresistances

$$R_{a,s}(H) \equiv \tfrac{1}{4}[R_{IV}(H) - R_{VI}(-H) + R_{IV}(-H) - R_{VI}(H)] \quad (1)$$

projects out that part of the elastic electron scattering due to the frozen spins and is called a spin fingerprint. Similarly choosing all plus signs in the above expression gives the OB symmetric and field symmetric $R_{s,s}(H)$, which is dominated by the scattering off nonmagnetic centers that accounts for at least 85% of the total scattering in these samples. What is the characteristic fluctuation field scale H_c of $R_{a,s}$? H_c is governed by two effects; one is orbital ($H_{\rm orb}$) and the other involves the Mn spins ($H_{\rm spin}$). Like quantum transport in nonmagnetic systems, the Aharonov-Bohm effect determines the characteristic orbital field $H_{\rm orb}$ from interfering Feynman paths: $H_{\rm orb}A \approx h/e$ (A the sample area projected perpendicular to H). At the same time, a change in H also distorts the spin configuration over the field scale $H_{\rm spin}$. H_c for $R_{a,s}$ is then given by $\min(H_{\rm orb}, H_{\rm spin})$. A detailed analysis [9] of $R_{a,s}$ in terms of invariants in spin and orbital space confirms this. Similarly $R_{s,s}$ reflects the preponderant nonmagnetic scattering and is experimentally found to be insensitive to spin reconfiguration induced by annealing and refreezing through T_g. Thus H_c for $R_{s,s}$ is given by $H_{\rm orb}$.

TABLE I. Fields are expressed in kG, and $\parallel$ and $\perp$ refer to the orientation of the field with respect to the wire. $H_{\rm spin}$, $H^{\parallel}_{\rm orb}$, and $H^{\perp}_{\rm orb}$ were determined from $R^{\parallel}_{a,s}$, $R^{\parallel}_{s,s}$, and $R^{\perp}_{s,s}$, respectively.

Concentration	Material parameter		Inferred quantity		
Atomic ppm Mn	T_g (K)	H_g	$H_{\rm spin}$	$H^{\parallel}_{\rm orb}$	$H^{\perp}_{\rm orb}$
320	0.5	2.7	1.3-2.1	14	2.0-2.5
1000	1.6	8.4	4.2-6.4	15	0.8-1.3

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Since H_{spin} is an interesting quantity in spin glasses [3], it is desirable to determine it directly from measurements of $R_{a,s}(H)$. We measured the correlation field H_c of $R_{a,s}$ while arranging $H_{orb} > H_{spin}$, so that $H_c = H_{spin}$. $H_{orb} \propto 1/A$ may be enhanced by decreasing the projected area A. For magnetotransport measurements with the applied field H perpendicular to the film plane $A_\perp = L_\phi(H)W$, while for H parallel to the wire $A_\parallel = Wd \ll A_\perp$. Figure 1 shows a spin fingerprint measured with H parallel to the wire (within $\sim 1°$). Such fingerprints are highly reproducible, having a correlation coefficient $C > 0.92$ between successive field sweeps at constant low temperature. The overall raw magnetoresistance $|\Delta R/R|$ was independent of field orientation and scaled appropriately with Mn concentration. Table I summarizes a series of experiments performed on wires of two concentrations. In both cases $H_c^\parallel$ for $R_{a,s}$ varied within the indicated range over the 0–90 kG scans, and was always smaller than $H_{orb}^\parallel$ as inferred from $R_{s,s}$. Thus $H_c^\parallel = H_{spin}$ is the field over which reversible changes in spin configuration occur. Apparently H_{spin} is less than or on the order of H_g for both alloys.

When an external field H is applied to a spin glass one expects that once H exceeds a few times H_g the magnetic ions are nearly all polarized. These ions are paramagnetic and follow $\mathbf{S}_i(-H) = -\mathbf{S}_i(H)$, so the amplitude of the fluctuations $\Delta g_{a,s} \equiv \Delta R_{a,s}/(R^2 e^2/h)$ should decay with increasing H. Figures 1 and 2 display spin fingerprints measured with two orientations of H. The amplitude of the $g_{a,s}$ fluctuations *does not vary* in either case even up to 90 kG $\approx 34 H_g$. Such experiments *seem* to indicate that the wires do not attain paramagnetic behavior even at surprisingly high fields.

Could some strong magnetic anisotropy, absent from bulk samples, account for this apparent high field spin rigidity? Let us first examine a surface anisotropy [10] K_s, which generates an equivalent anisotropy field $H_A = 4K_s v/\mu d$ (v the volume per Mn magnetic moment μ, d the film thickness or grain size). Since these high field effects are not observed in bulk millimeter sized samples, the bulk Dzyaloshinski-Moriya anisotropy (DMA) [11] sets an upper limit for surface effects: $K_s \ll 6 \times 10^{-1}$ erg/cm^2. Therefore for $d \simeq 400$ Å, $H_A \lesssim 6$ kG, which is much too weak to prevent polarization by 90 kG. Could an enhancement of the unidirectional DMA, perhaps arising from disorder effects [12] present in film samples, account for the rigidity effect? Any strong unidirectional anisotropy $H_A \gg H_g$, however, is inconsistent with the observed melting and refreezing of the low field spin fingerprints [4] near the bulk T_g. Because the $\Delta g_{a,s}$ fluctuation amplitude did not decay in large fields both in wires annealed for 5 min at 600 °C after deposition and in unannealed samples, growth induced anisotropies [13] are also precluded. Similarly the fact that the $\Delta g_{a,s}$ fluctuations did not diminish in samples with H applied both perpendicular and parallel to the film plane argues against stress induced magnetic anisotropies [14].

A remaining approach to the puzzle posed by Figs. 1 and 2 is to analyze the contributions to $\Delta g_{a,s}(H)$. This quantity contains both orbital and spin configuration distortion effects. But under a change of parallel field $\Delta H \approx H_{spin} < H_{orb}^\parallel$ starting at H, only the latter plays a role. The Lee-Altshuler theory [1] then predicts that the induced rms conductance change when a fraction $\xi(H, \Delta H)$ of the spins in an L_ϕ long section of wire rearranges is $\Delta G \simeq K L_\phi(H) \xi^{1/2}(H)$ (provided $\Delta G < e^2/h$), where K involves field independent material properties. A recent theoretical analysis [15] and our measurements of the length dependence of the conductance in magnetic

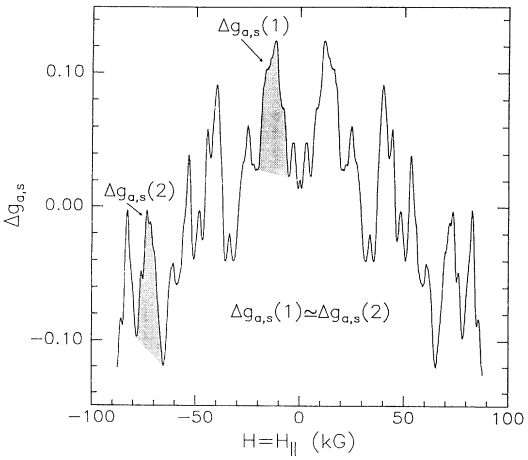

FIG. 1. A high field spin fingerprint $\Delta g_{a,s} \equiv \Delta R_{a,s}/[R^2(e^2/h)]$ measured on a 320 ppm 6500 Å×820 Å×370 Å *Cu*Mn wire with the field applied along the wire at 30 mK. The characteristic field scale for the fluctuations is due to distortion of the spin configuration. Notice that the amplitude of $\Delta g_{a,s}$ remains nearly unchanged.

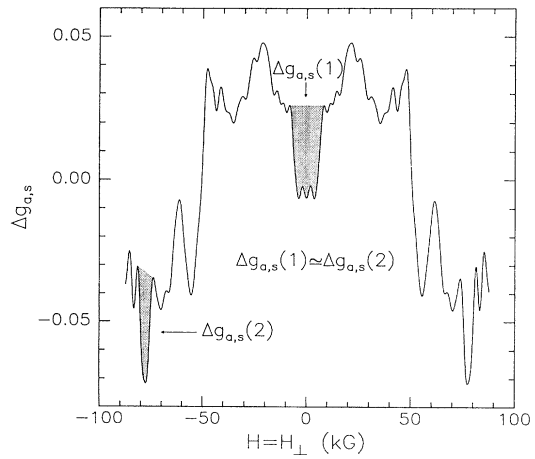

FIG. 2. A high field spin fingerprint measured on a wire of the same size and material but different resistance as in Fig. 1 with the field applied perpendicular to the wire. Again note that the fluctuation amplitude does not decay with field.

systems have shown that $\Delta R_{a,s}$ is nearly independent of wire length L for $L \gtrsim L_\phi$, so that $\Delta R_{a,s}(L) \simeq \Delta R_{a,s}(L_\phi) \simeq KR^2(L)(L_\phi/L)^2 L_\phi \xi^{1/2}$, where $R(L)$ is the wire resistance. Here L_ϕ varies with H because the dominant phase-breaking mechanism is spin-flip scattering off the Mn ions, which is gradually suppressed as the external field increases. $L_\phi(H)$ may be inferred from the correlation field $H_{\rm orb}^{\perp} = (h/e)/L_\phi W$ of the autocorrelation function of the fluctuations $\Delta R_{s,s}^{\perp}(H)$ in a 5–10 kG wide window centered on H. For the 1000 ppm material L_ϕ typically varies from 0.37 μm at low fields to 0.75 μm near 90 kG and is independent of the window size. Therefore the observed insensitivity of the observed $\Delta g_{a,s} = K(h/e^2)(L_\phi^3/L^2)\xi^{1/2}$ to field indicates nearly compensating field dependences for L_ϕ^3 and $\xi^{1/2}$.

We apply these ideas to learn about the distribution of exchange field magnitudes $P(h)$ (h the internal field strength in units of H_g). Each spin feels a local field which is the vector sum of the applied field H and an internal field generated by its interaction with the other spins. What do the spin fingerprints tell us about such distributions, and do a larger than naively expected fraction of the spins experience large molecular fields? Assume that the distribution $P(h)$ of internal fields is unaffected by H [16]. (We examine below the consequences of relaxing this constraint.) From the above experiments, we know the spins rearrange over changes in H on the order of $H_{\rm spin}$ with $H_{\rm spin} \lesssim H_g < H_{\rm orb}^{\parallel}$. Therefore imagine dividing the h axis of a plot of $P(h)$ vs h into $\Delta h = 1$ wide bins. Consider changing the parallel external field H from $H/H_g = n$ to $H/H_g = n+1$. Since $H_{\rm spin} < H_{\rm orb}^{\parallel}$, the fluctuations within each bin are primarily generated by spin rearrangement. Those spins feeling dimensionless *internal* fields $h > n+1$ will be held rigidly and will not contribute to the associated fluctuations $\Delta R_{a,s}$ connected with the changing applied field. Similarly those spins experiencing an internal field $h < n$ are already polarized before H is altered, and also do not produce a $\Delta R_{a,s}$. The spins that generate $\Delta R_{a,s}$ in this range of applied fields are just the fraction possessing internal fields with magnitudes from nH_g to $(n+1)H_g$, so that $P(H/H_g) = \xi(H) \simeq B[\Delta R_{a,s}(H)]^2 L_\phi^{-6}(H)$. Here $B = [L^2/KR^2(L)]^2$ is field independent since $|\Delta R(H)/R| \lesssim 3000$ ppm. Hence measuring $\Delta R_{a,s}^{\parallel}(H)$ and the field dependence of $L_\phi(H)$ directly yields $P(h)$, while B can be inferred from normalization.

Figure 3 illustrates such a distribution computed from measurements at 30 mK. The data represent a combination from a 320 and a 1000 ppm Mn device. We assumed the distribution $P(h = H/H_g)$ is the same for the two alloys, and joined the curves together at common values of h. The error bars refer to the range of variation occurring upon multiple melting and refreezing of the spin glass. An expected impurity averaged distribution for a randomly diluted alloy [17] is also depicted. The number of rearranging spins in each field bin ranges from 10^4 at $h \approx 1$–2 down to only 10^2 spins for $h > 20$.

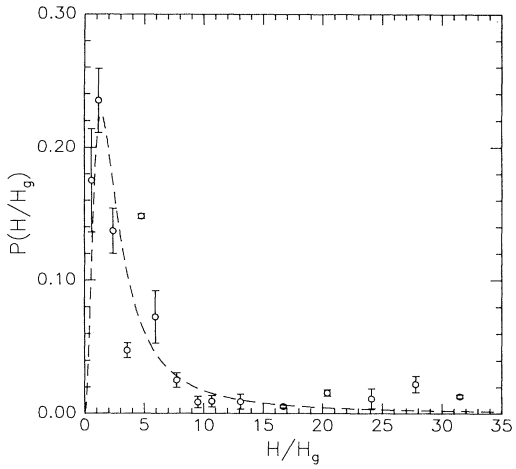

FIG. 3. The normalized distribution of internal fields in *Cu*Mn spin glass wires at 30 mK as computed from the spin fingerprints. The dashed curve is an expected ensemble averaged distribution for random Mn dilution: $P(h) = (4\sqrt{w}/\pi) \times [h^2/(w+h^2)^2]$, $h \equiv H/H_g$, adjusting the dimensionless width to $w = 2$.

The general features of such a $P(h)$ were robust to variations in experimental parameters. We studied five independently fabricated samples, each comprising wire sections of length L and resistance that ranged over an order of magnitude for the same concentration and found that $P(h)$ was largely unaltered. Also computing $P(h)$ only requires determining the relative field dependences of $\Delta R_{a,s}$ and L_ϕ rather than their absolute values due to the overall normalization condition. $P(h)$ was insensitive to temperature for $T/T_g \lesssim 0.6$–0.8. At higher temperatures the $\Delta R_{a,s}(H)$ traces became time dependent during the measurements, reflecting the molten spin system.

For $H/H_g > 20$ the experimental points in Fig. 3 lie systematically above $P(h)$ for random dilution. Including a "cavity" in the theoretical distribution [16] shifts some weight to larger fields. This enhances $P(h=20)$ by 8% and $P(h=35)$ by 23%, too small to account for the discrepancies. One may also use the known strength of the RKKY interaction in *Cu*Mn [17,18] to associate with each h a distance r corresponding to the separation between a pair of spins bonded to each other by a reduced field h. For $h = 20$–30 and $c = 320$ ppm, r ranges from 8 to 9 Å. For comparison, the Cu fcc unit cell size is 3.6 Å, and two Mn spins on the spatial diagonal experience $h \simeq 75$ for $c = 320$ ppm. Therefore for $h > 20$ the discreteness of the lattice plays a role not included in the dashed curve of Fig. 3. For $h = 25$–32 and $c = 320$ ppm, random dilution on the sites implies a probability $(2$–$3) \times 10^{-3}$ of having a pair bonded by a reduced field h. These are 4 to 7 times smaller than the experimental points, indicating that some short-range atomic ordering (correlated Mn arrangements) might be present. Neutron scattering experiments [19] have shown signs of atomic scale ferromagnetic regions containing a few spins

in bulk single crystal (5–20)% *Cu*Mn and also imply corrections of over 50% to the three nearest shell random Mn occupancies [20]. A direct comparison to this work, however, is complicated by the more than 50 times smaller concentrations used here.

The basic assumption of the analysis used to compute the internal field distribution $P(h)$ is that $P(h)$ is taken to be independent of H [16]. But as H increases, the internal field at a given spin changes as the spins to which it is coupled polarize. Thus H will induce probability flow from a bin $(h,h+1)$ into neighboring bins. One could expect that more flow will occur from bins with a higher starting weight $P_0 \equiv P(h,H=0)$. For P_0 decreasing with h, this will lead to a net shift of probability towards higher h as H increases. As a result, the measured $P(h)$ will have a longer tail than P_0.

In conclusion, a surprising robustness of spin fingerprints $\Delta g_{a,s}(H)$ was observed at large H/H_g and is attributed to their sensitivity to as few as 100 reorienting spins. A simple interpretation obtained for fields along the wires allows a separation of the spin and orbital contributions to the conductance fluctuations. This directly yields the characteristic spin configuration distortion field H_{spin} as well as the sample specific internal field distribution. The resulting distribution exhibits a relatively long tail that may be due to short-range magnetic order as found in bulk *Cu*Mn and/or to the influence of the external field on the internal fields. These results are examples of the versatility of this highly sensitive, configuration specific spin fingerprint technique.

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