Influence of an external magnetic field on the reentrant spin-glass transition temperature in a $Cr_{75}Fe_{25}$ alloy

S. M. Dubiel*
Max Planck Institut für Eisenforschung, D-4000 Düsseldorf, West Germany

K. H. Fischer, Ch. Sauer, and W. Zinn
Institut für Festkörperforschung, Kernforschungsanlage Jülich, D-5170 Jülich, West Germany
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The influence of an external magnetic field B_0 on the reentrant spin-glass transition temperature T_G of a $Cr_{75}Fe_{25}$ alloy has been investigated by means of ^{57}Fe Mössbauer-effect spectroscopy. A linear relationship in the (B_0,T) plane has been observed which is interpreted theoretically in terms of a mean-field-theory model including nonvanishing local spontaneous magnetization.

I. INTRODUCTION

According to the mean-field theory (MFT), spin-glass (SG) phase transitions occur also in the presence of an external magnetic field, B_0. In an Ising system, the transition in the (B_0,T) plane defines the so-called de Almeida–Thouless (AT) line[1] which is given by the following relationship:

$$\tau_c = Ah^{\theta}, \tag{1}$$

where $A=(\frac{3}{4})^{1/3}$, $\theta=\frac{2}{3}$, $\tau_c=1-T_G(B_0)/T_G(0)$, and $h=\mu_B g B_0/k_B T_G(0)$.

Equation (1) is valid for small h only, and $T_G(0)$ and $T_G(B_0)$ are the SG transition temperatures in zero field and in an external field, B_0, respectively. If the system is isotropic and consists of m-component spins, the transition in the (B_0,T) plane defines the Gabay-Toulouse (GT) line which in the approximation of small fields has the form[2]

$$\tau_c = Ch^{\theta} \tag{2}$$

where $C=(m^2+4m+2)/4(m+2)^2$ and $\theta=2$.

Although a large number of experiments based on different techniques have been performed for various systems[3–11] in order to determine the transition line(s) in the (B_0,T) plane, and hence to test the AT and GT predictions, a clear-cut picture has not yet emerged. A majority of the experimental data is in favor of the AT line, at least as far as the value of the exponent θ is concerned. However, the experimental data of the prefactor C are usually one order of magnitude larger than expected. On the other hand, there exists also some evidence which supports the GT line.[7]

The lack of perfect agreement between the experimental data and the two model predictions does not necessarily imply that the models are incorrect. The observed discrepancies can arise because the basic assumptions underlying those models often may not be fulfilled in real systems. For instance, Eq. (2) is strictly valid for the isotropic SG with a symmetric distribution of exchange interactions and with vanishing spontaneous magnetization $M_s=0$. Both, Eqs. (1) and (2) hold only for a model with infinite-range interactions. Real systems, however, may not fulfill these conditions. It has been shown[12] that random anisotropy plays an important role in the SG transitions, in particular in the presence of a magnetic field. For sufficiently strong random anisotropy, a crossover from the AT- to the GT-like behavior has been theoretically predicted[12] and also experimentally confirmed.[13] It is also known[2] for $B_0=0$ that the nonvanishing spontaneous magnetization affects the SG transition temperature, $T_G(0)$. The magnitude of the magnetization induced by B_0 is characteristic for a given system. Therefore, the experimental value of the parameter C is not universal, but depends on the given system.

In view of the above-mentioned complex situation, it seems to be important to provide further experimental data on the behavior of $T_G=f(B_0)$ for various real systems showing a SG transition. On the other hand, more realistic theoretical models are required, in order to improve the description of real SG systems. This is in particular true for reentrant transitions from a ferromagnetic to a spin-glass phase which is the subject of this paper.

Here, we present Mössbauer-effect data on the influence of an external magnetic field, B_0, on the transition temperature into the SG state for a $Cr_{75}Fe_{25}$ alloy. Furthermore, we try to explain the observed behavior in terms of the MF model of Sherrington and Kirkpatrick (SK).[14]

As we have previously shown by means of Mössbauer effect (ME) spectroscopy,[15] the sample of $Cr_{75}Fe_{25}$ exhibits a reentrant transition from a ferromagnetic (FM) into a SG state. In the ME experiment, the transition temperature can be obtained (i) from an anomaly in the average effective hf field, $\bar{B}$, when plotted versus the temperature, T, or (ii) from the intensity of the second and fifth lines of the ^{57}Fe ME spectrum, I_2 and I_5, respectively (i.e., those lines corresponding to the nuclear transitions with $\Delta m=0$). As demonstrated for the first time in Ref. 16 for Au-Fe, $\bar{B}$ shows a well-defined increase at a certain temperature, T_f, when plotted in the $(\bar{B},T)$ plane. This temperature was identified as the SG transition temperature, T_G. Figure 1 illustrates such a behavior for the $Cr_{75}Fe_{25}$

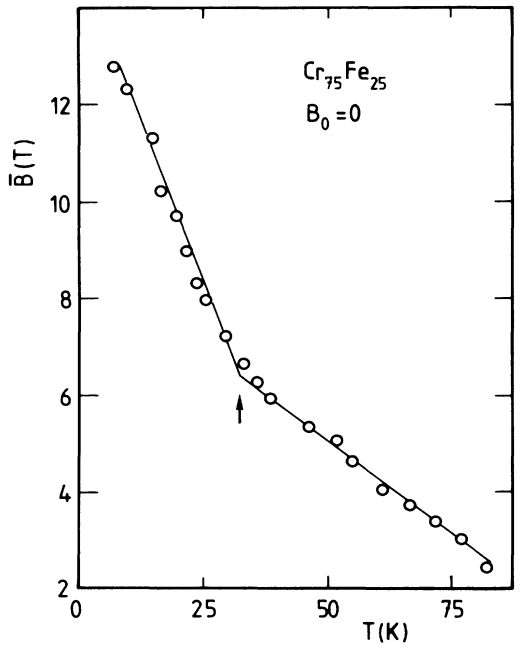

FIG. 1. Average effective hf field $\bar{B}$ vs the temperature, T, as obtained from zero-field ME measurements of $Cr_{75}Fe_{25}$ (Ref. 17). The arrow marks the transition temperature $T_G^{\bar{B}}$.

sample as measured in zero external field.[17] A value of $T_G(0)=30$ K can be readily deduced from this plot.

Concerning I_2 and I_5, these lines vanish if the propagation of the γ rays is parallel to the local magnetization vector μ. If, however, the γ rays propagate perpendicularly to μ, I_2 and I_5 exhibit their maximal values. In order to detect a FM→SG transition in that way, one has to align the magnetic moments in the FM phase by means of an external magnetic field, B_0, oriented parallel to the γ direction and to measure the Mössbauer spectra at different decreasing temperatures, T (field cooling). Entering the SG state at $T=T_G$, the moments will no longer be in parallel alignment and, consequently, I_2 and I_5 will now differ from zero.

The experimental results which are discussed in the present paper originate from new measurements in $B_0=3T$ as well as from our previous measurements: the zero-field data are taken from Ref. 17 and the $1T$ and $2T$ data from Ref. 15.

II. EXPERIMENTAL

A. Sample and Mössbauer effect measurements

We collected the ^{57}Fe Mössbauer effect spectra using the same $Cr_{75}Fe_{25}$ sample and the same setup as for the previous investigations.[15,17] We started the measurements, performed at different decreasing temperatures, by cooling the sample from $T\simeq 45$ K in a longitudinal magnetic field $B_0=3$ T. According to our previous study, the sample was at this starting temperature in the FM phase. Each spectrum was collected within a three-day run.

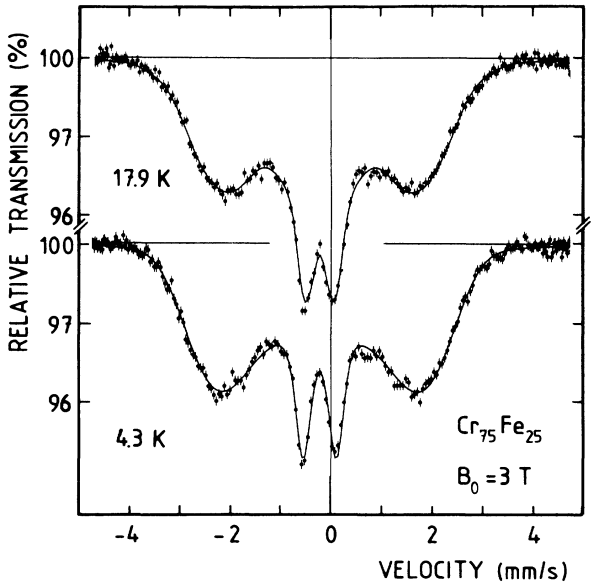

FIG. 2. Typical ^{57}Fe Mössbauer effect spectra of $Cr_{75}Fe_{25}$ as measured in an external magnetic field $B_0=3$ T at 17.9 K and 4.3 K. The solid line represents the fit using Window's method (Ref. 18).

B. Evaluation of the spectra

Two typical examples of the ^{57}Fe ME spectra are shown in Fig. 2. The quantitative evaluation of the spectra was based on the hf field-distribution method originally introduced by Window.[18] In order to derive the proper value of the x parameter ($x=I_{2,5}/I_{3,4}$) we used the same criterion as previously:[15] The standard deviation should be minimum by changing the x values stepwise as a free parameter, where we regarded only $P(B)$ curves without negative amplitudes [$P(B)$ is the distribution probability of the effective hf field at the ^{57}Fe nucleus, B]. This procedure has been already proved to work[15] as the $P(B)$ curves turned out to be very sensitive to the parameter x (typical error $\Delta x=\pm 0.05$). Figure 3 illustrates the $P(B)$ curves as obtained in this way from the ME spectra. The values of the average hf field were obtained from

$$\bar{B}=\int_0^{B_{\max}} P(B)B\,dB\ . \tag{3}$$

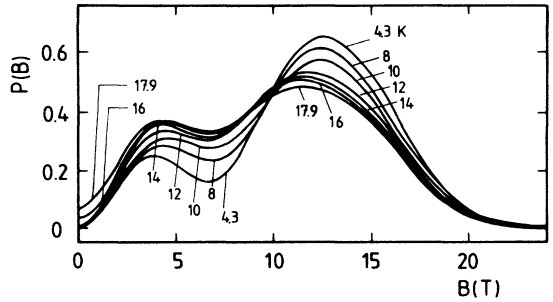

FIG. 3. Effective hyperfine field distribution curves $P(B)$ vs B as deduced from the ME spectra of $Cr_{75}Fe_{25}$ at different temperatures by means of Window's method (Ref. 18).

III. RESULTS

As mentioned in the Introduction, the transition temperature, T_G, can be deduced either from an anomaly in the behavior of $\bar{B}$ versus the temperature, T (valid for zero-field and in-field measurements) or from the x-versus-T dependence (valid only for in-field measurements). Figure 4 showing the temperature dependence of $\bar{B}$ demonstrates that the anomaly occurs here at $T_G^{\bar{B}}(3T)=12$ K. Also, as displayed in Fig. 5, a transition temperature of $T_G^x(3\text{ T})=14$ K can be inferred from the x-versus-T behavior. In this way, we have obtained two sets of data for the transition temperature: one, $T_G^{\bar{B}}$, from the anomaly in $\bar{B}$, and the other, T_G^x, deduced from the x-T plot. The question arises whether $T_G^{\bar{B}}$ and T_G^x measure the same transition. According to our present understanding of the issue, we think they do for the following reasons. $\bar{B}$ is proportional to the time-averaged value of the spin. In the FM phase ($T>T_G$) obviously only the longitudinal component (component parallel to $\mathbf{B}_0$) of the spin contributes to $\bar{B}$, while the transverse ones fluctuate much faster than the Larmor frequency. Hence its contribution to $\bar{B}$ vanishes. For $T<T_G$, i.e., in the SG state, the transverse components of the spin should be frozen, or, at least, fluctuate slower than the Larmor frequency, giving thereby an additional finite contribution to $\bar{B}$. Consequently, $\bar{B}$(SG) should be larger than $\bar{B}$(FM), as observed. In other words, the increase of $\bar{B}$ is directly related to the transversal component of the spin. Similarly, the x parameter remains zero as long as the γ rays and the local magnetization vectors μ are parallel to each other.

We see, however, from Fig. 5 that below a certain temperature, T_G^x, this is no longer the case. This may occur, either by a simple rotation of μ (without changing its amplitude) due to an increased anisotropy, or by the appearance of additional transverse quasistatic components of the spin. The two cases can be, however, simply distinguished by the behavior of $\bar{B}$. In the former case, $\bar{B}$ should not show any sudden increase at a certain temperature (T_G), which disagrees with our observations. We have, therefore, to conclude that it is the transversal component of the spin which causes that x becomes non-zero for $T\leq T_G$. In addition, as shown in our previous paper,[15] the x value measured at a given T depends strongly on the particular temperature and field cycling procedure. Especially, cooling in a larger B_0 produces larger x values. This further excludes the usual anisotropy to be responsible for $x\neq 0$ at $T\leq T_G$.

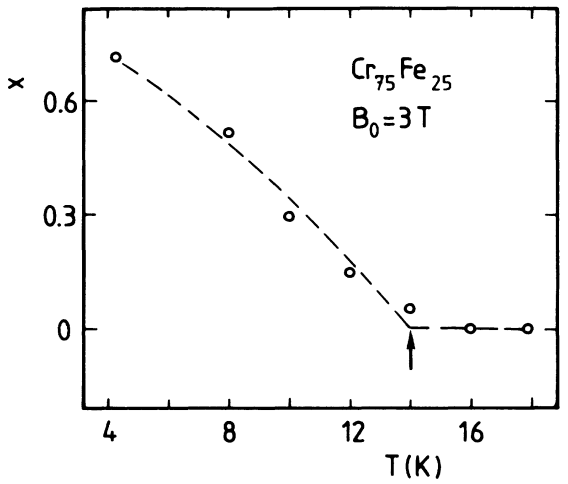

FIG. 5. Temperature dependence of the relative intensity x of the second and fifth resonance line for $Cr_{75}Fe_{25}$ measured in $B_0=3$ T. The arrow indicates the transition temperature T_G^x.

Following the arguments given above, we treat both $T_G^{\bar{B}}$ and T_G^x as temperatures marking the same transition. Thus, for the final discussion of the influence of B_0 on T_G, we take the mean value $T_G=\frac{1}{2}(T_G^x+T_G^{\bar{B}})$. Figure 6 shows T_G^x (diamonds), $T_G^{\bar{B}}$ (triangles), and the mean values T_G (circles) as a function of B_0. We see that the T_G^x data are close to $T_G^{\bar{B}}$. The existing difference between them can be considered as experimental errors involved in the determination of T_G by the present measurements. We also

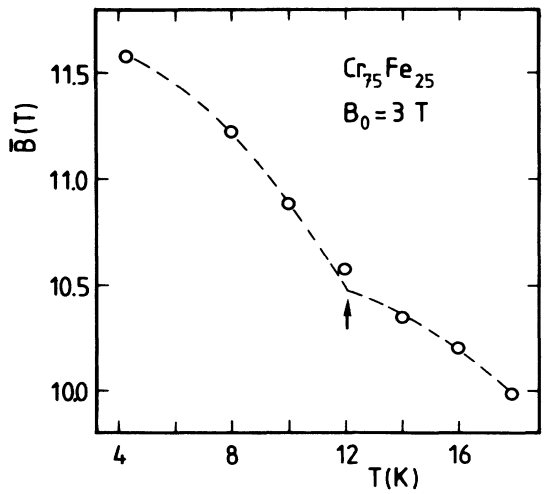

FIG. 4. Average effective hf field, $\bar{B}$ vs temperature T for $Cr_{75}Fe_{25}$ measured at $B_0=3$ T. The arrow marks the transition temperature $T_G^{\bar{B}}$.

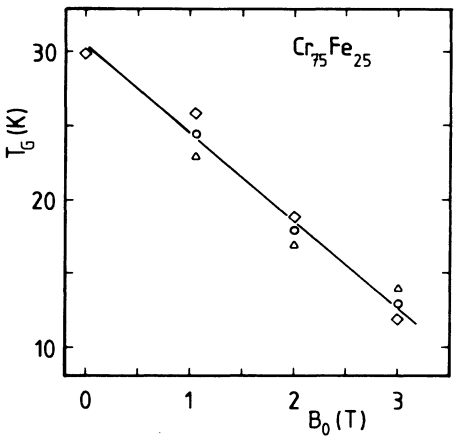

FIG. 6. Reentrant spin-glass transition temperature T_G as a function of the external magnetic field, B_0. Diamonds represent $T_G^{\bar{B}}$, triangles T_G^x, and circles $T_G=\frac{1}{2}(T_G^{\bar{B}}+T_G^x)$.

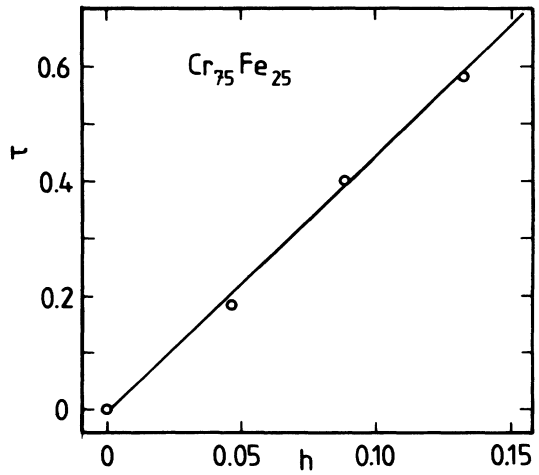

FIG. 7. Reduced reentrant spin-glass transition temperature $\tau=[1-T_G(B_0)/T_G(0)]$ vs $h=\mu_B g B_0/k_B T_G(0)$. The solid line is the best fit to the data: $\tau=[4.9(7)]h^{1.05(7)}$.

note that the T_G values decrease linearly with B_0.

For a quantitative comparison with the MFT predictions given in the next section, we present in Fig. 7 our data in the (τ,h) plane. The open circles represent our results and the solid line, being the best fit to the data, is given by

$$\tau_c(h)=1-\frac{T_G(B_0)}{T_G(0)}=(4.9\pm0.7)h^{1.05\pm0.07} . \quad (4)$$

We can conclude that in general our data are in agreement with the results of the MF model (see Sec. IV). The spin-glass temperature, T_G, decreases faster with B_0 than expected from the AT line and slower than expected from the GT line with a symmetric distribution of bonds. Also the prefactor is about five times larger than that derived from these models.

It is perhaps interesting to mention here that, based on magnetic viscosity measurements on $Cr_{84}Fe_{16}$ with $B_0 \leq 0.3$ T the authors of Ref. 6 found $\theta \simeq 0.4$. In order to compare this value with our result, one has to realize that (i) we investigated the FM→SG transition, i.e., the reentrant transition, while the sample investigated in Ref. 6 exhibited a simple paramagnetic SG transition (different sample compositions), and (ii) our external field was one order of magnitude larger. It is not clear whether the considerable difference in θ is due to the first or to the second reason. In the following section, we try to explain theoretically the linear correlation between T_G and B_0 which we have found.

IV. THEORY

The data discussed in Sec. III suggest the onset of quasistatic transverse local fields at $T_G(B_0)$ due to the canting of the spins out of the direction of the external field or remanent magnetization. Such an onset has been predicted by GT in MFT for isotropic vector spins in a field $\mathbf{B}_0$. The spin canting leads to a nonzero transverse order parameter $q_T=[\langle S_i^{\perp}\rangle^2]_J$. Here, $S_i^{\perp}$ is the component of a spin $\mathbf{S}_i$ at site i perpendicular to the external field $\mathbf{B}_0$, $\langle \cdots \rangle$ means thermal averaging and $[\cdots]_J$ the average over a Gaussian distribution of bonds

$$P(J_{ij})=\left[\frac{N}{2\pi J^2}\right]^{1/2}\exp\left[-\frac{N}{2J^2}(J_{ij}-J_0/N)^2\right] . \quad (5)$$

The GT line is defined by the condition $q_T\to 0$, and (2) is obtained[2] for a *symmetric* distribution (5) with $J_0=0$. This result is based on the SK model[14] for classical vector spins with infinite-range interactions and with the Hamiltonian

$$H=-\tfrac{1}{2}\sum_{i,j}J_{ij}\mathbf{S}_i\cdot\mathbf{S}_j-g\mu_B\mathbf{B}_0\cdot\sum_i\mathbf{S}_i . \quad (6)$$

For $J_0>J$, the model (5) and (6) exhibits a double transition.[2,14] With decreasing temperature, the system becomes first ferromagnetic (with nonzero spontaneous magnetization $\mathbf{M}_s$) and then a spin glass. The spin-glass transition $T_G(B_0)$ defined by $q_T\to 0$ marks also the onset of irreversibility effects such as the extremely slow decay of the magnetization or of spin correlations. However, the spontaneous magnetization M_s does not vanish at $T_G(0)$, in contrast to the conventional definition of a spin-glass state. This has been shown for Ising spins[14] but holds also for vector spins (see below).

The onset of irreversibility along the GT line[2] is relatively weak, as shown for a symmetric distribution (5) with $J_0=0$.[19,20] One has in vector spin glasses at lower fields and temperatures a second line which marks the crossover from weak to strong irreversibility and has the field-temperature dependence of the AT line (1). One expects such a line to exist also in reentrant systems below the line $T_G(B_0)$. The spontaneous magnetization decreases in the spin-glass region of reentrant systems with decreasing temperature, but seems to be finite at all temperatures.

Random anisotropy leads to a crossover from Heisenberg to Ising-type spins.[12] For $B_0\neq 0$, one has a coupling of longitudinal and transverse modes, and both q_L and q_T are nonzero already above the transition line $T_G(B_0)$. The latter no longer is defined by $q_T\neq 0$ or spin canting but only by the onset of irreversibility or an instability of the solution with a single vector order parameter $\mathbf{q}(T,B_0)$.[21] The data on $Cr_{75}Fe_{25}$ presented in Sec. III indicate that in this system the anisotropy is sufficiently small and can be ignored, i.e., one has a clear onset of spin canting for $B_0\neq 0$.

Surprisingly, the line $T_G(B_0)$ for reentrant isotropic vector systems has not yet been calculated. We will show that the ferromagnetic state is strongly distorted in reentrant systems. The field dependence of $T_G(B_0)$ is influenced by the spontaneous magnetization $\mathbf{M}_s$ and differs from both the AT and GT lines (1) and (2). The magnetization $\mathbf{M}(T,B_0)$ and the transverse and longitudinal spin-glass order parameters $q_T(T,B_0)$ and $q_L(T,B_0)$ have to be determined self-consistently. We restrict ourselves to the discussion of fields and temperatures on and above the line $T_G(B_0)$ where no replica symmetry breaking[21] is needed.

One of the authors and others derived[19] the partition function for vector spin glasses from (5) and (6) with $\mathbf{B}_0=(B_0,0,0)$, $\mathbf{M}=(M,0,0)$ (in units of $g\mu_B$), and $J=k_B T_G(0)$. We have

$$Z=m^{1/2}(2\pi)^{(m-1)/2}\int_{-m^{1/2}}^{m^{1/2}}dS_1(m-S_1^2)^{(m-3)/4}\,|a|_{m-1}^{(3-m)/2}I_{(m-3)/2}(|a|_{m-1}(m-S_1^2)^{1/2})\exp(a_1S_1+bS_1^2)\ , \tag{7}$$

with $h_{\text{eff}}=Jh+J_0M$, $q_1=q_L$, $q_2=\cdots q_m=q_T$, and

$$a_\alpha=\beta(h_{\text{eff}}\delta_{\alpha,1}+z_\alpha q_\alpha^{1/2})\ ,\quad \alpha=1,\ldots,m \tag{8}$$

$$b=\tfrac{1}{2}\beta^2(q_T-q_L+mX)\ , \tag{9}$$

$$|a|_{m-1}=(a_2^2+\cdots+a_m^2)^{1/2}\ . \tag{10}$$

Here, $I_m(z)$ are modified Bessel functions of the first kind. One derives from (7) the magnetization M, the spin-glass parameters q_α and the quadrupolar parameter X (Ref. 19)

$$M=\left[Z^{-1}\frac{\partial Z}{\partial a_1}\right]_J\equiv\int_{-\infty}^{\infty}\prod_\alpha(dz_\alpha/2\pi)Z^{-1}\frac{\partial Z}{\partial a_1}\exp\left[-\tfrac{1}{2}\sum_\alpha z_\alpha^2\right]\ , \tag{11}$$

$$q_L=\left[\left[Z^{-1}\frac{\partial Z}{\partial a_1}\right]^2\right]_J\ ,\quad q_T=\left[\left[Z^{-1}\frac{\partial Z}{\partial a_2}\right]^2\right]_J\ , \tag{12}$$

$$X=\left[1-Z^{-1}\frac{\partial^2 Z}{\partial a_2^2}\right]_J\ . \tag{13}$$

We expand (7) to order h_{eff}^2, q_L^2, and q_T. For Heisenberg spins with $m=3$ this leads to the coupled self-consistent equations:

$$q_L[1-(\beta J)^2]=(\beta h_{\text{eff}})^2-\tfrac{2}{5}(\beta J)^4(5q_L^2-6q_LX)-\tfrac{16}{5}(\beta J)^2(\beta h_{\text{eff}})^2q_L\ , \tag{14}$$

$$q_T=(\beta J)^2q_T\left[1-\tfrac{2}{5}\beta^2(h_{\text{eff}}^2+3J^2X)+\beta^4\frac{3\times53}{4\times5\times5\times7}q_L^2\right]\ , \tag{15}$$

$$X=\tfrac{1}{5}\beta^2(h_{\text{eff}}^2+3J^2X)-\tfrac{1}{5}(\beta J)^4q_L^2\ . \tag{16}$$

The Eqs. (14) to (16) lead to

$$q_L=\tfrac{1}{4}((\beta J)^2-1+\{[(\beta J)^2-1]^2+8(\beta h_{\text{eff}})^2\}^{1/2}) \tag{17}$$

and[19]

$$X=\tfrac{1}{4}(\beta h_{\text{eff}})^2\ ,\quad [T\approx T_G(B_0)]\ . \tag{18}$$

The Eqs. (14) to (18) hold in a region where the spontaneous magnetization is small or zero. Near the Curie temperature $T_C=J_0/k_B$ (see below) and for $J_0/J=T_C/T_G\gg1$ and $B_0=0$ one has from (17), $q_L=(\beta J_0)^2M^2\approx M^2$: All spin-glass properties vanish and one has an ideal ferromagnet. In what follows, we consider the neighborhood of the critical line $T\approx T_G(B_0)$ or $\tau^2\ll2(\beta h_{\text{eff}})^2$ where $\tau=1-T/T_G(0)$. One has in this limit from (17)

$$q_L=2^{-1/2}\beta|h_{\text{eff}}|+\tfrac{1}{2}\tau+\cdots\ . \tag{19}$$

The opposite limit $1\gg\tau^2\gg2(\beta h_{\text{eff}})^2$, with

$$q_L=\tau+\tfrac{5}{2}\tau^2+(\beta h_{\text{eff}})^2/2\tau\quad(\tau>0)\ , \tag{20a}$$

$$q_L=\tfrac{5}{2}\tau^2+(\beta h_{\text{eff}})^2/2|\tau|\quad(\tau<0)\ , \tag{20b}$$

is important near T_C, $B_0\to0$, and would be important for an AT-like line. Equations (15), (18), and (19) lead[2] to the generalization of the GT line (2):

$$\tau_c(h)=C(\beta h_{\text{eff}})^2=C(h+MJ_0/J)^2\ , \tag{21}$$

with $k_BT\approx k_BT_G(0)=J$.

The magnetization (11) is to order h_{eff}^3

$$M=\beta h_{\text{eff}}[1-(\beta J)^2(q_L-\tfrac{6}{5}X)-\tfrac{1}{5}(\beta h_{\text{eff}})^2]\ . \tag{22}$$

For $J=0$, one has the usual Landau expansion which leads to the Curie temperature $T_C=J_0/k_B$, to the spontaneous magnetization

$$M_s=5^{1/2}[(T_C-T)/T_C]^{1/2}$$

near T_C and to the induced magnetization $M\sim B_0^{1/\delta}$ with $\delta=3$ at T_C.

Near the triple point $J=J_0$ or $T_G(0)=T_C$, this behavior is strongly modified. Near T_C, one has with (20b) the quadrupolar parameter, Eq. (16), $X=\frac{1}{2}(\beta h_{\text{eff}})^2$ which leads to

$$M_s=(5^{1/2}/2)[(T_C-T)/T_C)]^{1/2}$$

or one-half of the magnetization for $J=0$. Near the critical line (21) holds $\tau^2\ll2(\beta h_{\text{eff}})^2$ and with (19),

$$M_s=2^{1/2}(T_C-T)T_G(0)/T_C^2\ , \tag{23}$$

i.e., the spontaneous magnetization varies *linearly* with T_C-T: the tendency of the interactions (5) to align the spins randomly as in a spin glass leads to a fairly distorted ferromagnetic state already above the spin-glass temperature $T_G(0)$. Such a "frustrated" ferromagnetic state has been investigated[22] in considerable detail by neutron scattering on the reentrant system $Eu_{0.54}Sr_{0.46}S$. Formally, this linear temperature-dependence arises from the linear field dependence of q_L, Eq. (19), together with (22).

Equation (23) leads to the critical line (21) for $h=0$ and $T_C \approx T_G(0)$

$$\tau_c(0)=2C(T_C-T_G)^2/T_G^2=2C(J_0-J)^2/J^2 , \quad (24)$$

which agrees[2] with the result of GT.

The magnetization M for a finite field B_0 is also easily calculated. One has near the triple point $J=J_0$ from (22)

$$M+h=\tfrac{1}{2}M_s+(\tfrac{1}{4}M_s^2+2^{1/2}h)^{1/2} . \quad (25)$$

For $h=0$, this leads to $M=M_s$, at the triple point with $M_s=0$ to $M=2^{1/4}h^{1/2}+O(h)$, and for $h \ll M_s$ to $M=2^{1/2}h/M_s$. The longitudinal susceptibility diverges at the triple point $J=J_0$

$$\chi=\frac{dM}{dh}\sim h^{1/\delta} , \quad \delta=2 , \quad (26)$$

in contrast to an ideal ferromagnet with $\delta=3$ at the Curie temperature and in contrast to a spin glass with the exponent $\delta=2$ for the *nonlinear* susceptibility $\chi_{\rm nl}\sim h^{2/\delta}$ at $T_G(0)$. All these critical exponents hold in MFT.

The critical line $\tau_c(h)$ is now easily calculated. One has for $h \ll M_s$

$$\tau_c(h)=\tau_c(0)+2\times 2^{1/2}Ch \quad (27)$$

and for the triple point (with $M_s=0$), $\tau_c(h)=2^{1/2}Ch$. Here, the constant C is defined in (2) and $\tau_C(0)$ in (24). The critical line $\tau_c(h)$ now varies *linearly* with the external field $B_0\sim h$, in contrast to the AT line (1) or the GT line (2). In all cases, one has a *decrease* in the critical temperature with increasing field. The difference between (27) and the GT line (2) can be traced back to the induced and/or spontaneous magnetization $M(T,h)$ which either is irrelevant or zero for a symmetric distribution (5) with $J_0=0$. For $J_0=0$, Eq. (21) leads back to the result (2) of GT.

V. CONCLUSIONS

The linear field dependence derived in (27) agrees well with the result (4) from ^{57}Fe ME spectra on $Cr_{75}Fe_{25}$. The double transition of this system is well established.[15,17,23–25] However, in comparing the experimental data with the theory presented in Sec. IV, several aspects and limitations should be noted.

(1) The results presented in Sec. IV are based on the SK model[14] which describes a MFT. The MFT predicts only part of the properties of real spin glasses correctly,[26] and there are even indications that three-dimensional isotropic Heisenberg spin glasses with short-range interactions and $J_0=0$ have no finite-temperature phase transition,[27] in contrast to Ising spin glasses[28] or systems with random anisotropy.[29] The critical exponents observed in some spin glasses (presumably systems with additional anisotropy) differ considerably from those predicted in MFT.

(2) The model (6) with a random distribution of bonds represents also for short-range interactions a strong simplification of true spin glasses. For instance, it is not obvious how the parameter $J_0/J=T_C/T_G(0)$ is related to the Fe concentration of Cr-Fe. The theory says nothing about the concentration dependence of the Curie temperature T_C or freezing temperature T_G. The phase diagrams of systems such as Au-Fe, Eu-Sr-S, and Cr-Fe are partly determined by percolation, an aspect which is completely missing in our model.

(3) The expansion (22) holds only for small magnetizations. This assumption is correct near the Curie temperature T_C and any ratio of $T_C/T_G(0)$. However, it holds at $T_G(B_0)$ only for $T_C-T_G(0) \ll 1$ or near the triple point $J=J_0$. Fortunately, our basic result $\tau_c(h)=\tau_c(0)+\alpha h$ with some constant α holds also more generally since it follows directly from (21), together with linear response relation $M=M_s+\chi h$. Our result is then restricted either to the region of linear response ($J_0>J$) or to the triple point $J=J_0$.

(4) Some experiments seem to indicate the vanishing of the spontaneous magnetization M_s below a certain temperature T_G'.[26] At present, it is unclear whether T_G' agrees or disagrees with the temperature $T_G(B_0)$. The Mössbauer data discussed in this paper yield no information about the *global* magnetization $M(T)$ but only about local properties. Magnetization measurements on $Cr_{75}Fe_{25}$ are hampered by a homogeneous distribution of magnetization,[30] or do not go to sufficiently low temperatures[30] or fields.[31] The MFT presented in Sec. IV predicts $M_s \neq 0$ at the critical line $T_G(B_0)$, apart from the region $J_0 \leq J$. A test of this prediction would be extremely interesting.

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*Permanent address: Academy of Mining and Metallurgy, PL-30 059 Krakow, Poland.

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