Mean-field theory of freezing temperatures of two-component metallic spin glasses

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A mean-field approach which emphasizes spin-spin pair correlations is used to calculate the freezing temperature T_g of two-component metallic spin glasses. We calculate the variation in T_g as a function of the ratio of concentrations of the two magnetic species. Finite-temperature effects and intrinsic *sd* scattering of electrons play important roles in establishing the trends exhibited by T_g. A clear physical interpretation is given which provides insight into the behavior of T_g in real, two-component spin glasses. The theory satisfactorily reproduces experimental results without the use of any adjustable parameters. We also predict how T_g will behave if the concentrations of the two species are varied while the *ratio* c_2/c_1 of the concentrations is held fixed.

I. INTRODUCTION

Despite the inherent complexity of spin glasses, significant understanding has been obtained through the application of seemingly "simple" theoretical methods. An excellent example is the use of mean-field theory to explain the cusp in the magnetic susceptibility χ of spin glasses (see for example Refs. 1–4). Intense investigations by many researchers have provided significant advances in the understanding of spin glasses. Nevertheless, numerous intriguing questions remain to be solved.[5]

The primary purpose of this paper is to contribute to the understanding of the physical processes which determine the properties of spin glasses by addressing the question of the behavior of the freezing temperatures T_g of two-component metallic spin glasses. We show that an improved understanding of the trends in T_g is obtained by applying mean-field theory.

The behavior of T_g in two-component spin glasses is a very complex question. The decision to use mean-field theory in an attempt to describe the trends in T_g in two-component systems is motivated by successes of the type described above. More specifically, recent investigations of the trends in T_g (Refs. 6–8) and χ (Ref. 9) of metallic spin glasses through the use of mean-field theory has proven to be rather successful. Since we will apply the methods of these studies, we begin by briefly outlining the key aspects of these works.

A new approach[6-8] was recently proposed for calculating the impurity concentration dependence of the freezing temperature of a single-component metallic spin glass such as AgMn or CuMn. The essential features of the approach are (i) to include the intrinsic *sd* scattering of conduction electrons from the magnetic ions (spins) in the spin glass, (ii) to explicitly retain, and to reveal the importance of, finite temperature effects, and (iii) to emphasize the role of spin-spin correlations and in doing so to derive a new expression for T_g. The expression obtained is

$$T_g \propto \frac{1}{N} \sum_{i,j}{}' ([J_{ij}^2]_{av})^{1/2} , \tag{1}$$

where J_{ij} is the effective indirect exchange coupling between a pair (i,j) of spins in the spin glass, $[\cdots]_{av}$ denotes an average over all possible configurations of magnetic ions and nonmagnetic ions (if present), N is the total number of spins, and the prime on the sum means to exclude $j = i$.

Equation (1) gives excellent agreement with several experimental results. For example, excellent agreement between theory and experiment is obtained for the dependence of T_g on the concentration c of magnetic ions in metallic spin glasses such as AgMn or CuMn. The theory also provides a very good fit to experimental data exhibiting the dependence of T_g on the concentration c_i of nonmagnetic impurities in spin glasses with c held fixed (e.g., AgMn with Mn concentration fixed but varying concentrations of Sb). These and other comparisons of the theory with experimental results are discussed more fully in Refs. 6–8. The new theoretical approach also explains the cusp in the isothermal dc magnetic susceptibility observed experimentally at $T = T_g$ (Ref. 9).

The primary purpose of this paper is to extend the alternative theoretical approach to calculate $T_g(c_1, c_2)$ in two-component metallic spin glasses in which the concentration of one species (type "1") is c_1 and the concentration of the other species (type "2") is c_2. We restrict our investigations to two-component metallic spin glasses in which the intrinsic *sd* scattering is the dominant scattering mechanism. We exclude consideration of various other possibilities including, for example, multicomponent spin glasses in which spin-orbit scattering must be taken into account.[10] We take c_1 to be nonzero, finite, and fixed, and we treat c_2 as variable. We focus attention on the range $0 < x \equiv c_2/c_1 \lesssim 3$ so as to be in accord with the range typically investigated experimentally.[11] We place emphasis on the very important features (i)–(iii) cited above.

II. A MEAN-FIELD THEORY FOR TWO-COMPONENT SPIN GLASSES

Mean-field theory provides the following expressions for $|\langle \mathbf{S}_i^{(1)} \rangle|$ and $|\langle \mathbf{S}_i^{(2)} \rangle|$ when the magnetic ions are in a

particular configuration:

$$|\langle S_i^{(1)} \rangle| = S_1 B_{S_1}(S_1[g_i^{(1)} + g_i^{(12)}]/kT) , \qquad (2)$$

and

$$|\langle S_{i'}^{(2)} \rangle| = S_2 B_{S_2}(S_2[g_{i'}^{(2)} + g_{i'}^{(12)}]/kT) , \qquad (3)$$

where $\langle \cdots \rangle$ denotes the thermal expectation value at temperature T,

$$g_i^{(1)} = \sum_j {}'[J_{ij}^{(1)}|\langle S_j^{(1)} \rangle| \cos\phi_{ij}] , \qquad (4)$$

$$g_i^{(12)} = \sum_{j'} [J_{ij'}^{(12)}|\langle S_{j'}^{(2)} \rangle| \cos\phi_{ij'}] , \qquad (5)$$

with similar expressions for $g_{i'}^{(2)}$ and $g_{i'}^{(12)}$, and where B_S is the well known Brillouin function, the prime on the sum means to exclude $j=i$, ϕ_{ij} is the angle between $\langle S_i^{(1)} \rangle$ and $\langle S_j^{(1)} \rangle$, and other angles are defined similarly.

In what follows we will take T to be just less than T_g. Consequently, all of the $|\langle S_i^{(1)} \rangle|$ and $|\langle S_{i'}^{(2)} \rangle|$ are small, and we may expand the Brillouin functions in Eqs. (2) and (3). Expanding Eq. (2) to lowest order, multiplying both sides by $(1/N_1)|\langle S_i^{(1)} \rangle|$, and summing over all spins "i" of type "1" gives

$$m_1^2(T) \equiv \frac{1}{N_1} \sum_i |\langle S_i^{(1)} \rangle|^2 = \frac{S_1(S_1+1)}{3k_B T} G_1(T) , \qquad (6)$$

where

$$G_1(T) \equiv \frac{1}{N_1} \sum_i \sum_j {}'[J_{ij}^{(1)}\langle S_i^{(1)} \rangle \cdot \langle S_j^{(1)} \rangle]$$
$$+ \frac{1}{N_1} \sum_i \sum_{j'} [J_{ij'}^{(12)}\langle S_i^{(1)} \rangle \cdot \langle S_{j'}^{(2)} \rangle] . \qquad (7)$$

Similarly,

$$m_2^2(T) \equiv \frac{1}{N_2} \sum_{i'} |\langle S_{i'}^{(2)} \rangle|^2 = \frac{S_2(S_2+1)}{3k_B T} G_2(T) , \qquad (8)$$

where

$$G_2(T) \equiv \frac{1}{N_2} \sum_{i'} \sum_{j'} {}'[J_{i'j'}^{(2)}\langle S_{i'}^{(2)} \rangle \cdot \langle S_{j'}^{(2)} \rangle]$$
$$+ \frac{1}{N_2} \sum_i \sum_{j'} [J_{ij'}^{(12)}\langle S_i^{(1)} \rangle \cdot \langle S_{j'}^{(2)} \rangle] ; \qquad (9)$$

N_1 and N_2 are, respectively, the numbers of spins of type "1" and of type "2".

The importance of spin-spin pair correlations is captured in Eqs. (6)–(9). Consider for example the case of no type "2" impurities ($c_2=0, c_1>0$). If $T=T_g^-$, then $m_1^2>0$, and we immediately have from Eq. (6) that $G_1(T_g^-)>0$. Using Eq. (7), $G_1(T_g^-)>0$ shows that the spins are pairwise correlated: Each spin pair has a *tendency* to assume a relative orientation such that $[J_{ij}^{(1)}\langle S_i^{(1)} \rangle \cdot \langle S_j^{(1)} \rangle]>0$. Indeed, completely random spin directions would give $G_1(T)=0$ even for $T<T_g$ and there would be no spin-glass state. Equation (1) for T_g was derived by exploiting these spin-pair correlations (see Refs. 6–8). These correlations play a major role in the

derivation of $T_g(c_1,c_2)$ given in the remainder of this paper.

Some assumptions and approximations are required at this stage. Making use of the spin-spin pair correlations, we approximate $G_1(T)$ and $G_2(T)$ by

$$G_1(T) \propto m_1^2(T)H_1(T) + \frac{c_2}{c_1}m_1(T)m_2(T)H_{12}(T) , \qquad (10)$$

and

$$G_2(T) \propto m_2^2(T)H_2(T) + m_1(T)m_2(T)H_{12}(T) , \qquad (11)$$

where

$$H_1(T) = \left[\frac{1}{N_1} \sum_i \sum_j {}'|J_{ij}^{(1)}|\right]_T , \qquad (12)$$

$$H_2(T) = \left[\frac{1}{N_2} \sum_{i'} \sum_{j'} {}'|J_{i'j'}^{(2)}|\right]_T , \qquad (13)$$

and

$$H_{12}(T) = \left[\frac{1}{N_2} \sum_{j'} \sum_i |J_{ij'}^{(12)}|\right]_T . \qquad (14)$$

For simplicity we will assume that the constants of proportionality in Eqs. (10) and (11) are equal; this is equivalent to taking the overall frustration between pairs of spins to be independent of the type ("1" or "2") of each spin. Combining Eqs. (6), (8), (10), and (11), we obtain two equations for T_g:

$$T_g \propto \frac{S_1(S_1+1)}{3k_B}\left[H_1(T_g) + \frac{c_2}{c_1}qH_{12}(T_g)\right] , \qquad (15)$$

$$T_g \propto \frac{S_2(S_2+1)}{3k_B}\left[H_2(T_g) + \frac{1}{q}H_{12}(T_g)\right] , \qquad (16)$$

where the constants of proportionality are again equal and

$$q \equiv m_2(T_g^-)/m_1(T_g^-) . \qquad (17)$$

Equations (15) and (16) will be regarded as the two equations which determine the values of the two quantities T_g and q; as discussed below, all other quantities will turn out to be "knowns."

The coupling J_{ij} has been calculated previously.[6-8] Taking into account the intrinsic sd scattering of electrons from magnetic ions in the spin glass and explicitly retaining the finite-T dependence gives

$$|J_{ij}| = h(R_{ij})K_0(R_{ij}) , \qquad (18)$$

where $K_0(R)$ is the coupling strength between two spins in an otherwise pure system, $[K_0(R)=2mk_F j_{sd}^2/(2\pi R)^3$ for $R \gg k_F^{-1}]$, j_{sd} is the sd exchange constant which enters the portion of the Hamiltonian that describes the sd scattering, and a simplified version of $h(R)$ was shown to be

$$h(R) = \begin{cases} 1, & \lambda_{nn} < R < \Lambda_{sd}, & (19) \\ \gamma, & \Lambda_{sd} < R < \Lambda_T, & (20) \\ 0, & R > \Lambda_T, & (21) \end{cases}$$

where λ_{nn} is the mean nearest-neighbor distance, Λ_T is the well known finite interaction range at temperature T,

$$\Lambda_T = \left[\frac{\lambda T_F}{3\pi k_F T_g}\right]^{1/2} , \tag{22}$$

λ is the electron mean free path, Λ_{sd} is the length scale which arises as a consequence of the intrinsic sd scattering, and γ is a constant which has the approximate value 0.6062 (Ref. 8).

Equations (18)–(21) give the coupling strength between a typical pair of spins separated by distance R. We may use these equations to obtain a clear physical picture of the pairwise spin couplings. This picture will prove to be extremely useful for interpreting the trends in $T_g(c_1,c_2)$ which we will derive below. We see that we may regard all spin pairs having separation $R < \Lambda_{sd}$ as being coupled with strength $K_0(R)$, all pairs having $R > \Lambda_T$ as being completely uncoupled, and all pairs having $\Lambda_{sd} < R < \Lambda_T$ as having coupling strength $\gamma K_0(R)$. Since $0 < \gamma < 1$, we see that the sd scattering has the effect of giving a *reduction* in the spin-spin coupling strength at $R = \Lambda_{sd}$.

It was also shown previously[8] that, for concentrations of order of or less than a few atomic percent, Λ_{sd} was only slightly larger than λ. In the remainder of this paper, we will take $\Lambda_{sd} = \lambda$.

We use Eqs. (18)–(21) for $|J_{ij}^{(1)}|$, $|J_{i'j'}^{(2)}|$, and $|J_{ij'}^{(12)}|$ simply by making the proper choice for λ_{nn} in each case. For $H_1(T)$, $H_2(T)$, and $H_{12}(T)$ we use, respectively, $\lambda_{nn} = \lambda_{nn}^{(1)}$, $\lambda_{nn} = \lambda_{nn}^{(2)}$, and $\lambda_{nn} = \lambda_{nn}^{(12)}$; $\lambda_{nn}^{(1)}$ is the mean nearest-neighbor distance between type "1" pairs, $\lambda_{nn}^{(2)}$ between type "2" pairs, and $\lambda_{nn}^{(12)}$ is most easily defined by describing how the sum $H_{12}(T)$ will be evaluated. We choose to evaluate $H_{12}(T)$ as follows. For a given spin j' of type "2" we perform the sum over all spins "i" of type "1," then average the result over all spins of type "2." We may therefore represent $H_{12}(T)$ by an integral over distances $R_{j'i}$ from spins of type "2" to spins of type "1." (See Ref. 8 for details.) The lower limit of this integral will then be the mean distance from a type "2" spin to a type "1" spin. This choice is appropriate because we have chosen c_1 as fixed and we regard c_2 as increasing from 0; consequently, when $c_2 \ll c_1$, type "2" spins will be surrounded by type "1" spins and are themselves much farther apart. As such, assuming for simplicity that the spins are positioned randomly, we see that the lower limit $\lambda_{nn}^{(12)}$ is simply $\lambda_{nn}^{(1)}$. Using Eqs. (18)–(21) for the coupling constants and evaluating the sums in Eqs. (12)–(14) gives, after some algebra,

$$t_g(x) = [1 + j_{12}^2 x q(x)] f(x) , \tag{23}$$

and

$$t_g(x) = s\left[j_2^2 x + \frac{j_{12}^2}{q(x)}\right] f(x) + \frac{1}{3} s j_2^2 \alpha_1 x \ln x , \tag{24}$$

where

$$f(x) \equiv 1 - \alpha_1\left[\left(1 - \frac{\gamma}{2}\right)\ln(1+bx) + \frac{\gamma}{2}\ln[t_g(x)]\right] , \tag{25}$$

$$t_g(x) \equiv T_g / T_g^{(1)} , \tag{26}$$

$$x \equiv c_2/c_1 , \tag{27}$$

$$s \equiv [S_2(S_2+1)]/[S_1(S_1+1)] , \tag{28}$$

$$j_2 \equiv [j_{sd}^{(2)}]/[j_{sd}^{(1)}] , \tag{29}$$

$$j_{12} \equiv [j_{sd}^{(12)}]/[j_{sd}^{(1)}] , \tag{30}$$

$$\alpha_1 \equiv \left[\ln(\lambda_1 c_1^{1/3}) + \frac{\gamma}{2}\ln\left[\frac{T_F}{3\pi k_F \lambda_1 T_g^{(1)}}\right]\right]^{-1} , \tag{31}$$

$$b \equiv B_1/B_2 , \tag{32}$$

with the physical meaning of $q(x)$ given by Eq. (17), $T_g^{(1)} \equiv T_g(c_1, c_2=0)$ is the freezing temperature when $c_2 = 0$ and $c_1 > 0$, and $\lambda_1 = B_1/c_1$ and $\lambda_2 = B_2/c_2$ are, respectively, the electron mean free paths when only type "1" or type "2" spins are present.

In deriving Eqs. (23) and (24) we have assumed for simplicity that the spin positions are random, and have therefore taken $\lambda_{nn}^{(1)} = c_1^{-1/3}$ and $\lambda_{nn}^{(2)} = c_2^{-1/3}$. We have also implicitly assumed that $\lambda_{nn}^{(1)} < \lambda$, $\lambda_{nn}^{(2)} < \lambda$, and $\lambda < \Lambda_{T_g}$. These inequalities restrict the range of x to

$$x_{min} < x < x_{max} \tag{33}$$

where x_{min} is given by

$$x_{min}^{-1/3}(1 + bx_{min}) = \lambda_1 c_1^{1/3} \tag{34}$$

and

$$x_{max} = (\lambda_1 c_1^{1/3} - 1)/b ; \tag{35}$$

note that x_{min} and x_{max} are functions of c_1. The restriction is not severe; for example, the values $b = 1$, $\lambda_1 = 100$ Å, and $c_1 = 0.001$ Å^{-3} give $0.001 \lesssim x < 9$, which is certainly suitable for our purposes.[12]

Note that all of the following quantities can be obtained from experimental data and/or theoretical formulas: S_1, S_2, j_{12}, j_2, b, λ_1, T_F, k_F, and $T_g^{(1)}$ (details are given below). Also recall that we treat c_1 as given and fixed and c_2 as variable. Consequently, Eqs. (23) and (24) give $T_g(x)$ and $q(x)$ as functions of x with the various quantities appearing in the two equations having values characteristic of the particular spin glass. We emphasize that, in using these two equations to compare theory with experimental data for a given two-component spin glass, there will be *no free parameters* to adjust to give a fit (see below).

We solve Eqs. (23) and (24) by using Eq. (24) to express $q(x)$ in terms of $t_g(x)$ and x, and inserting the result into Eq. (23) to give a transcendental equation for $t_g(x)$:

$$q(x) = \frac{s j_{12}^2 f(x)}{t_g(x) - s j_2^2 x [f(x) + (1/3)\alpha_1 \ln x]} \tag{36}$$

and

$$t_g(x) = f(x) + \frac{s j_{12}^4 x [f(x)]^2}{t_g(x) - s j_2^2 x [f(x) + (1/3)\alpha_1 \ln x]} . \tag{37}$$

An equivalent but simpler expression for $q(x)$ is obtained

from Eq. (23):

$$q(x) = \frac{t_g(x) - f(x)}{j_{12}^2 x f(x)} , \qquad (38)$$

where we now regard $t_g(x)$ as determined by Eq. (37).

Equations (37) and (38) are the principal results of this paper. They are readily solved numerically for given values of the various quantities in them.

III. PHYSICAL DISCUSSION

In order to reveal the fundamental physical features captured by the theory, we show in Figs. 1–3 how $t_g(x)$ and $q(x)$ vary with the quantities j_2 and α_1. Note that this heuristic investigation is to be distinguished from the comparison of theory and experiment of the type presented below. In Figs. 1–3 we take $j_{12}^2 = j_2$ (viz. we take $[j_{sd}^{(12)}]^2 = [j_{sd}^{(1)}][j_{sd}^{(2)}]$).

Figure 1 shows $t_g(x)$ vs x for various values of j_2. The three principal features of the curves are (1) $t_g(x)$ increases with j_2 at all x, (2) for small enough j_2, $t_g(x)$ initially decreases with x, reaches a minimum, then increases with x, and (3) for large enough j_2, $t_g(x)$ is monotonic increasing with x.

These three features are readily interpreted in terms of competition between two physical effects. (I) As type "2" spins are added, the nearest-neighbor spin-spin distances decrease, and the total concentration of moments increases; both tend to elevate T_g. (II) The addition of type "2" spins, however, results in a decrease in both of the important length scales $\Lambda_{sd} = \lambda$ and Λ_{T_g}; consequently, there is a reduction in the spin-spin coupling strengths of
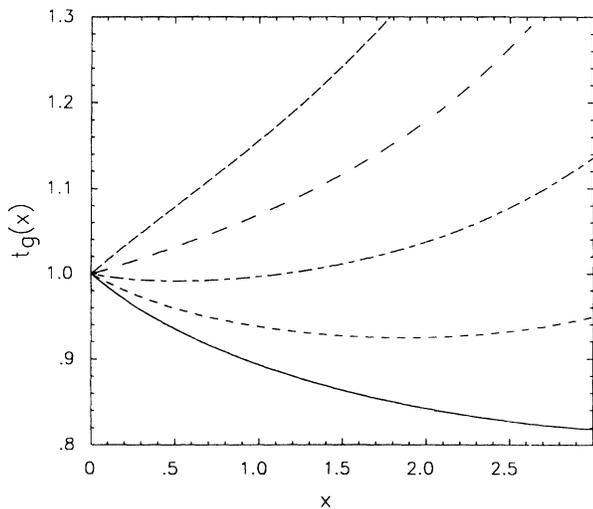

FIG. 1. The freezing temperature $t_g(x) \equiv T_g(c_1,c_2)/T_g(c_1,c_2=0)$ as a function of $x \equiv c_2/c_1$ in a two-component metallic spin glass with magnetic ion concentrations c_1 and c_2. The curves shown are for five values of the ratio j_2 of sd coupling strengths between the two types of spins, $j_2 \equiv j_{sd}^{(2)}/j_{sd}^{(1)}$. From bottom curve to top, the values of j_2 are 0.25, 0.35, 0.45, 0.55, and 0.65. The values of the other quantities which determine $t_g(x)$ are $s = 1$, $b = 1$, $\gamma/2 = 0.3031$, and $\alpha_1 = 0.3554$; see text for details.
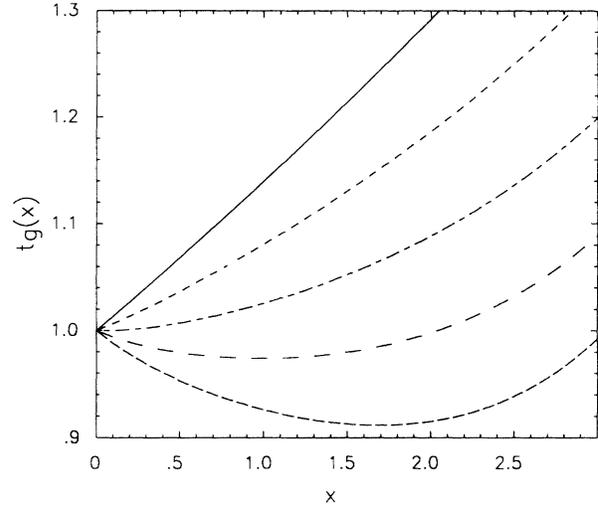

FIG. 2. As in Fig. 1 but with $j_2 = 1$, $s = 1$, $b = 1$, $\gamma/2 = 0.3031$, and with values of α_1 from top curve to bottom as follows: 0.1, 0.2, 0.3, 0.4, 0.5.

some pairs (due to the decrease in Λ_{sd}), and a decoupling of other pairs (due to Λ_{T_g} decreasing), with the result that T_g tends to drop. [Recall the physical interpretation given just after Eq. (22).]

As a specific example, the feature (1) in Fig. 1 results because an increase in j_2 means enhanced spin-spin pair interactions and an elevation in T_g; in the case of feature (2), the effect of reduced length scales dominates initially, causing T_g to fall, but eventually gives way to pair-coupling effects, at which stage T_g rises; in (3) the length scale effects are always less significant than pair-coupling effects and T_g is monotonic increasing. Note that the
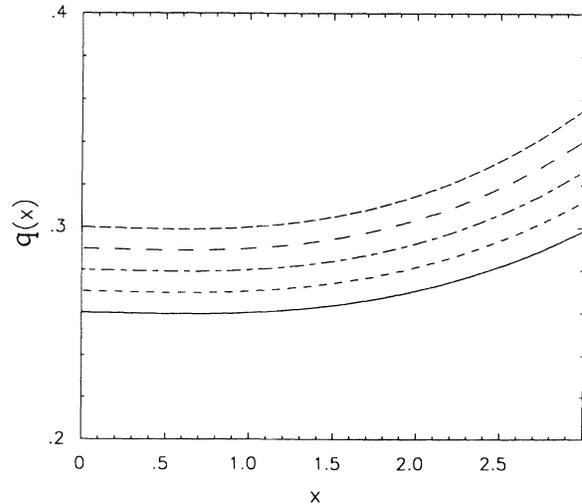

FIG. 3. The ratio of magnitudes of spin $q(x) \equiv m_2(T_g^-)/m_1(T_g^-)$ as a function of $x \equiv c_2/c_1$ in a two-component metallic spin glass with magnetic ion concentrations c_1, c_2. As in Fig. 1, the five curves, from bottom to top, have the following values of j_2: 0.52, 0.54, 0.56, 0.58, 0.60. The other quantities have the values $\alpha_1 = 0.3$, $\gamma/2 = 0.3$, $b = 1$, and $s = 0.5$. See text for details.

effects due to reductions in the length scales show up in all the $t_g(x)$ curves in Fig. 1. Also, note that the choices $s = 1$ and $b = 1$ have been made to isolate the effect of variation in j_2, while the value $\alpha_1 = 0.3554$ is simply a realistic choice, corresponding arbitrarily to $Ag_{0.974}Mn_{0.026}$ (viz. Mn is type "1").

The competition between the two physical effects explains the primary features in Fig. 2 as well, where α_1 is varied. In Fig. 2, the trend is *opposite* to that of Fig. 1: $t_g(x)$ *decreases* with increasing α_1. If s or γ is varied the trends are like those in Fig. 1, while the character of curves for different values of b are as in Fig. 2. (Note that even though γ has a fixed value, we are free to vary γ for heuristic purposes.) In all cases the physical explanations are similar to those given above for changes in j_2.

We wish to point out that the results shown in Fig. 2 represent a prediction of our theory. Specifically, an increase in α_1 may be thought of as due to an increase in c_1 [see Eq. (31) and note that $\ln(\lambda_1 T_g^{(1)})$] varies weakly compared to $\ln(\lambda_1 c_1^{1/3}) = \ln(B_1 c_1^{-2/3})$]. Consequently, our theory predicts that $t_g(x)$ will decrease with c_1 in a fashion similar to that of Fig. 2. This is understood physically as follows. If c_1 is varied with all other quantities held fixed (including x), only two length scales vary independently, say λ and $\lambda_{nn}^{(1)}$. Since λ decreases more rapidly than $\lambda_{nn}^{(1)}$, the reduction in $t_g(x)$ due to reduction in spin-spin coupling strength dominates over the tendency for $t_g(x)$ to increase due to diminished nearest-neighbor distance. The mathematical explanation is most easily revealed by using Eq. (37) to show that

$$t_g'(0^+) \equiv \left[\frac{\partial t_g(x)}{\partial x} \right]_{x \to 0^+} = \frac{sj_{12}^4 - \alpha_1(1 - \gamma/2)b}{1 + \alpha_1 \gamma/2} , \quad (39)$$

which in turn gives

$$\frac{\partial [t_g'(0^+)]}{\partial \alpha_1} = - \frac{[sj_{12}^4 \gamma/2 + (1 - \gamma/2)b]}{(1 + \alpha_1 \gamma/2)^2} ; \quad (40)$$

this means that $\partial [t_g'(0^+)]/\partial c_1$ is always negative. In terms of experimental behavior of a given two-component metallic spin glass in which sd scattering dominates, our prediction is as follows. In different samples having the same ratio c_2/c_1, as c_1 increases the $t_g(x)$ curves for the different samples will vary in the sense from the top curve down in Fig. 2.

In Fig. 3 we show the variation in $q(x)$ with j_2. Note that the values $q(0^+)$ are given by

$$q(0^+) = sj_2 \quad (41)$$

[see Eq. (36) and recall that $j_{12}^2 = j_2$ in Figs. 1–3]. The physical explanation for the features in Fig. 3 closely parallels that for Fig. 1.

We also point out that the limiting case where c_1 is held fixed and $c_2 \to 0$ has been examined[12] to ensure that we recover the results reported previously[8] for single-component metallic spin glasses.

IV. COMPARISON OF THEORY WITH EXPERIMENT

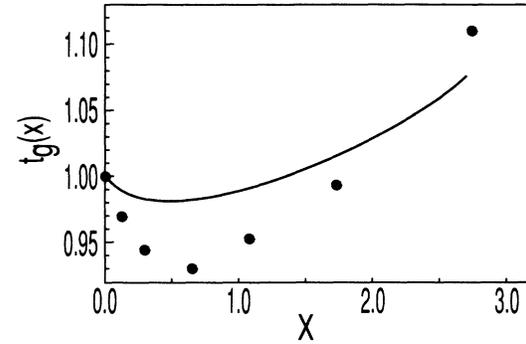We briefly compare our theory with experimental results. Two types of behavior have been reported in two-

FIG. 4. The freezing temperature $t_g(x)$ in the two-component spin glass AuCrFe. The Cr concentration is fixed at 2.3 at. %, while the concentration of Fe varies. The solid dots are data presented in Ref. 11; the curve is the theoretical result, Eq. (37). See text for details.

component metallic spin glasses.[11] One is that $T_g(c_1, c_2)$ decreases initially as c_2 increases with c_1 held fixed, reaches a minimum, then increases with c_2. This is shown in Fig. 4 for AuCrFe where the concentration of Cr is fixed while that of Fe increases. Our theory (solid curve) provides a satisfactory account of the experimental behavior. The other type of behavior is that T_g increases monotonically with c_2. Figure 5 compares our theory with experimental data for AuFeMn with the concentration of Fe fixed and that of Mn varying. The agreement between theory and experiment is rather good, and it appears from Figs. 4 and 5 that the theory has captured the essential physics: The physical explanation underlying these behaviors is the competition between the two opposing physical effects described above. (A brief physical discussion was also given in Ref. 8.) We stress that the theoretical curves in Figs. 4 and 5 contain *no adjustable parameters*.

The curves in Figs. 4 and 5 were obtained using Eqs. (37) and (25), with values of (sj_2^2), α_1, and b obtained from the literature in the following manner. Values for (sj_2^2) were obtained by using Eqs. (15) and (16) to show that
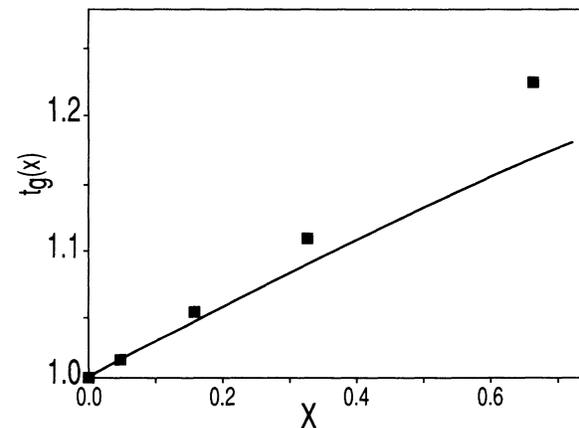

FIG. 5. As in Fig. 4 but for AuFeMn with the concentration of Fe fixed at 3.0 at. % and the Mn concentration varying. The solid squares are data given in Ref. 11, while the curve is the theoretical result, Eq. (37).

$$(sj_2^2) = \frac{c_1 \alpha_2(c_2) T_g^{(2)}(c_2)}{c_2 \alpha_1(c_1) T_g^{(1)}(c_1)} , \qquad (42)$$

where $\alpha_2(c_2)$ and $T_g^{(2)}(c_2)$ are, respectively, the analogs of $\alpha_1(c_1)$ and $T_g^{(1)}(c_1)$ when only species "2" is present with concentration c_2. Values of $T_g^{(1)}(c_1)$ and $T_g^{(2)}(c_2)$ are readily found in the literature.[11,13] Values of $\alpha_1(c_1)$ and $\alpha_2(c_2)$ were obtained using Eq. (31) and calculating $\lambda_1(c_1)$ or $\lambda_2(c_2)$ using a free-electron picture[14] and having values for the resistivities $\rho_1(c_1)$ or $\rho_2(c_2)$ for the specific material.[15] These resistivity values also provide values for b. In Fig. 4, $(sj_2^2)=0.50$, $\alpha_1=0.48$, and $b=1.76$; in Fig. 5, $(sj_2^2)=0.40$, $\alpha_1=0.55$, and $b=0.31$. (Note that we again take $j_{12}^2 = j_2$ in Figs. 4 and 5.)

Finally, we note that ideas along the lines of those presented in this paper were mentioned briefly by Vier and Schultz.[11] Henley extended their presentation by considering a mean-field theory of multicomponent spin glasses.[16] However, a basic assumption was that the range of interaction was infinite. Although mean-field theories with infinite interaction ranges can give reasonable results for transition temperatures, we have seen that the finite range of interaction Λ_T plays a prominent role in the physical processes which establish the behavior in spin glasses, and we therefore regard our approach as preferable.

V. SUMMARY

In summary, we have presented a derivation of a mean-field theory for the spin-glass freezing temperature in a two-component metallic spin glass, and we have illustrated the basic underlying physical ideas. We have also shown that our theory satisfactorily accounts for experimental results and provides a physical understanding of these results. Our calculations included the sd scattering of conduction electrons by the magnetic ions and finite temperature effects. Emphasis was placed on the role of spin-spin pair correlations. We also predict that the initial slope of $t_g(x)$ for a given two-component metallic spin glass will decrease as c_1 increases.

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