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To cite this version:

V. Thanh Ngo, D. Tien Hoang, Hung T. Diep, I. A. Campbell. Ising Spin Glass on the FCC lattice: Phase Diagram and Relaxation. 2013. hal-00860682v1

HAL Id: hal-00860682
https://hal.archives-ouvertes.fr/hal-00860682v1

Submitted on 10 Sep 2013 (v1), last revised 3 Mar 2014 (v2)
Ising Spin Glass on the FCC lattice: Phase Diagram and Relaxation

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We study the phase transition in a face-centered-cubic antiferromagnet with Ising spins as a function of the concentration $p$ of ferromagnetic bonds randomly introduced into the system. Such a model describes the spin-glass phase at strong bond disorder. Using the standard Monte Carlo simulation and the powerful Wang-Landau flat-histogram method, we carry out in this work intensive simulations in the whole range of $p$. We show that the first-order transition disappears with a tiny amount of ferromagnetic bonds, namely $p \approx 0.01$, in agreement with theories and simulations on other 3D models. The antiferromagnetic long-range order is also destroyed with a very small $p$ ($\geq 5\%$). With increasing $p$, the system changes into a spin glass and then to ferromagnetic phase when $p > 0.65$. The phase diagram in the space $(T, p)$ shows a dissymmetry, unlike the case $\pm J$ Ising spin glass on the simple cubic lattice. We calculate the relaxation time around the spin-glass transition temperature and we show that the spin autocorrelation follows a stretched exponential relation law where the factor $b$ is equal to $\simeq 1/3$ at the transition as suggested by the percolation-based theory. This value is in agreement with experiments performed on various spin glasses and with Monte Carlo simulations on different SG models.

PACS numbers: 75.10.Nr ; 75.10.-b ; 75.40.Mg ; 64.60.-i

I. INTRODUCTION

Spin glasses\textsuperscript{1–3} (SG) have been a subject of intensive investigations for more than four decades. The main difficulties in such systems come from the combination of the frustration\textsuperscript{4} and the bond disorder.

One of the most interesting questions is the law of the relaxation time in disordered systems. There is a large number of theories yielding various results since the first pioneer work by R. Kohlrausch which introduced phenomenologically in 1847 the so-called stretched exponential relaxation (SER) law. The reader is referred to the review by Phillips\textsuperscript{5} for numerous examples of systems, not limited to SG, with theoretical and experimental results. We are interested in this paper in SG; there have been several theoretical and numerical works dealing with the SER. Ogielski\textsuperscript{6} has studied an Ising spin model of SG on a simple cubic lattice with a random nearest-neighbor $\pm J$ bond distribution by using intensive Monte Carlo (MC) simulations. Among his numerous results, he found that dynamic correlations follow the SER above the spin-glass transition temperature $T_g$ and a power law below $T_g$ for short and long times. The SER is written as

$$q(t) = \langle S(0)S(t) \rangle = Ct^{-x} \exp(-\omega t^b) \quad (1)$$

where $S(t)$ is an Ising spin at the time $t$, $\langle ... \rangle$ indicates the thermal average, $< ... >=$ denotes an average over disordered samples, $C$ a constant, $x$ and $b$ are temperature-dependent coefficients. Ogielski found that $b$ is equal to $\simeq 1/3$ at $T_g$ and it increases with increasing temperature ($T$). Below $T_g$, he found that $q(t)$ follows the algebraic decay $Ct^{-x}$ where $x$ does not have the same temperature dependence above and below $T_g$. de Dominicis \textit{et al.}\textsuperscript{7} found by a calculation in the random free energy landscape that the relaxation to equilibrium follows a SER which depends on the choice of the transition probability between valleys in the free energy space. There has been a number of other theories showing a SER behavior among which we can mention the trap model\textsuperscript{8}, the hierarchical model\textsuperscript{9}, and the model of random walkers on dilute hypercubes of high dimensions\textsuperscript{10–13}. In the last model, Almeida \textit{et al.}\textsuperscript{13} found, by a calculation using a random walk on the dilute hypercube, a SER near the percolation transition which is similar to the topology of spin configurations near the transition in Ising SG. They found that $b$ is equal to 1/3 at the percolation threshold $p_c$ and it increases with increasing $p$. A recent MC simulation of the Sherrington-Kirkpatrick Ising SG model at and above $T_g$ shows a long-time SER with $b$ equal to $\simeq 1/3$ at $T_g$\textsuperscript{14}. Experimentally, let us mention a few data which show a SER with precise values of $b$. Early works on canonical SG $\text{AgMn}\textsuperscript{15,16}$ found a SER below $T_g$ with $b = 1 - n$ where $n \approx 0.62$ for $T/T_g$ from 0.5 to 0.8. In an experiment on $\text{Cu}_{0.5}\text{Co}_{0.5}\text{Cl}_2\text{FeCl}_3$ graphite bi-intercalation compound, Suzuki and Suzuki\textsuperscript{17} observed a $b$ value nearly equal to 0.3 at $T_g$. Malinowski \textit{et al.}\textsuperscript{18} found in Ni-doped $\text{La}_{0.5}\text{Sr}_{0.5}\text{CoO}_4$ a best fit of their time dependence of the relaxation with a SER.
Another point which motivates our present work is the effect of quenched bond randomness in systems having a first-order transition. Note that in a pure system with a second-order transition, the effect of a small amount of bond disorder is understood according to the Harris criterion\(^\text{19}\). In systems with a first-order transition, this question has been examined since 1989 by a number of important investigations mostly in two dimensions (2D)\(^\text{20–23}\). These works have shown that a small quenched bond disorder suffices to make disappear the latent heat in a pure 2D system having a first-order transition. In particular, Hui and Berker\(^\text{22}\) have demonstrated this by looking at the Blume-Emery-Griffiths (BEG) spin-1 model in 2D, with a renormalization group (RG) argument. Cardy and Jacobsen\(^\text{23}\) have illustrated this with the \(q = 8\) Potts model by the use of finite-size scaling and conformal invariance. There is however a contradiction between different researchers concerning the universality of the resulting second-order transition upon introducing a small disorder: some authors\(^\text{22–26}\) claimed that this should belong to a new random fixed point for \(q > 2\), while others favored the 2D Ising universality class\(^\text{27}\). In 3D, Falicov and Berker\(^\text{28–29}\) have shown by a RG calculation that the tricritical point in the pure BEG model is replaced, when bond randomness is introduced, by a segment of second-order transition bounded at one end by a multicritical point and at the other end by a new random tricritical point. Interestingly, they have shown that there exists a threshold of randomness beyond which the random tricritical point goes down to \(T = 0\). In a recent paper, Malakis \textit{et al.}\(^\text{29}\) have carried out a MC study on the 3D BEG model on a simple cubic lattice. They showed that the tricritical point moves in the direction to reduce the first-order line in the phase diagram and the critical behavior of the ex first-order transition line belong to the 3D random Ising model. In addition, Fernandez \textit{et al.}\(^\text{30}\) have also shown by microcanonical MC simulations that the 3D site-dilute 4- and 8-state Potts models confirm the conjecture of Cardy and Jacobsen\(^\text{23}\) in 3D. For quantum spin systems, Greenblatt \textit{et al.}\(^\text{31}\) showed rigorously the disappearance of first-order transition with randomness in 2D and in dimension \(d \leq 4\), just as for the classical case\(^\text{21}\). Note that in the case where the pure model has a second-order transition such as when \(q = 4\), Domany and Wiseman\(^\text{32}\) have shown that a random quenched bond disorder changes the second-order \(q = 4\) Potts universality into a 2D pure Ising one.

The present paper addresses the two points mentioned above: (i) to verify the conjecture of Cardy and Jacobsen\(^\text{23}\) on the disappearance of the latent heat in the case of a 3D first-order transition when a quenched bond disorder is introduced (see Fig. 1 of their paper), and to see if the result found for the 3D BEG model\(^\text{28–29}\) is applied to other models such as the one studied in the present paper (ii) to study dynamic correlations of the spin-glass phase created at a strong bond disorder of our model to verify that the results of Ogields are model-independent. To carry out these purposes, we consider the pure Ising face-centered cubic (FCC) antiferromagnet with nearest-neighbor (NN) interaction. This model is fully frustrated since the lattice is composed of equilateral triangles with antiferromagnetic interaction\(^4\). It shows a very strong first-order transition\(^\text{23–26}\).

The paper is organized as follows. Section II is devoted to a description of our model and the simulation method. Results are shown and discussed in section III. Concluding remarks are given in section IV.

\section{Model and Wang-Landau Algorithm} \label{sec:method}

We consider the FCC lattice with Ising spins of magnitude \(S = 1\). The Hamiltonian is given by

\begin{equation}
\mathcal{H} = - \sum_{\langle i,j \rangle} J_{ij} S_i S_j
\end{equation}

where \(S_i\) is the Ising spin at the lattice site \(i\), \(\sum_{\langle i,j \rangle}\) is made over the NN spin pairs \(S_i\) and \(S_j\) with interaction \(J_{ij}\). Hereafter we suppose that \(J_{ij} = -J (J > 0)\) for antiferromagnetic bonds and \(J_{ij} = J\) for ferromagnetic bonds. The lattice is composed of \(L^3\) FCC-cells each with four spins (\(L\) being the number of cells in each direction). The four sublattices are shown in Fig. 1. The total number of spins of the lattice is \(N = 4L^3\). We use the periodic boundary conditions. The number of NN bonds per spin is 12, thus the total number of bond of the system is \(N_b = 12 \times N/2 = 24L^3\). Hereafter, the energy is measured in the unit of \(|J| = 1\) and \(T\) in the unit of \(|J|/k_B = 1\)

We start with the pure antiferromagnetic state. For this case, it is well-known that the bulk FCC antiferromagnet with Ising spins shows a very strong first-order transition\(^\text{33–36}\). The phase transition in the Heisenberg model also shows a first-order character\(^\text{37,38}\). Other similar frustrated antiferromagnets such as the HCP antiferromagnet show the same behavior\(^\text{39,40}\).

To create a spin glass, we introduce a number of ferromagnetic bonds \(N_b^F\) in a random manner into the pure antiferromagnetic state. The ferromagnetic bond concentration \(p\) is thus \(p = N_b^F / N_b\). The phase transi-
tion behavior of the system depends on the parameter \( p \) (0 \( \leq p \leq 1 \)), i.e. it changes from pure antiferromagnetic phase at \( p = 0 \) to pure ferromagnetic phase at \( p = 1 \), passing through the spin-glass phase.

In order to investigate the nature of the phase transition and the relaxation behavior at various \( p \), we use the standard MC method and the Wang-Landau technique of simulation. Wang and Landau\(^4\) recently proposed a MC algorithm for classical statistical models which allowed to study systems with diﬃcultly accessed microscopic states. In particular, it allows us to detect with eﬃciency weak ﬁrst-order transitions\(^4\).\(^2\).\(^3\). The algorithm uses a random walk in the energy space to get an accurate estimate for the density of states (DOS) \( g(E) \) which is defined as the number of spin conﬁgurations for any given \( E \). This method is based on the fact that a ﬂat energy histogram \( H(E) \) is produced if the probability for the transition to a state of energy \( E \) is proportional to \( g(E)^{-1} \). We summarize how this algorithm is implied here. At the beginning of the simulation, the density of states \( g(E) \) is unknown so all densities are set to unity, \( g(E) = 1 \). We begin a random walk in energy space \( (E) \) by choosing a site randomly and ﬂipping its spin with a transition probability
\[
p(E \rightarrow E') = \min \left[ g(E)/g(E'), 1 \right],
\]
where \( E \) is the energy of the current state and \( E' \) is the energy of the proposed new state. Each time an energy level \( E \) is visited, the DOS is modiﬁed by a modiﬁcation factor \( f > 0 \) whether the spin is ﬂipped or not, i.e. \( g(E) \rightarrow g(E)f \). At the beginning of the random walk, the modiﬁcation factor \( f \) can be as large as \( e^1 \simeq 2.7182818 \). A histogram \( H(E) \) records the number of times a state of energy \( E \) is visited. Each time the energy histogram satisﬁes a certain “ﬂatness” criterion, the histogram \( H(E) \) is then reset to zero, and the modiﬁcation factor is reduced, typically to the square root of the previous factor, to produce a ﬁner estimate of \( g(E) \). The reduction process of the modiﬁcation factor \( f \) is repeated several times until a ﬁnal value \( f_{\text{final}} \) which is close enough to one. The histogram is considered as ﬂat if
\[
H(E) \geq x\% \cdot \langle H(E) \rangle
\]
for all energies, where \( x\% \) is chosen between 90\% and 95\% and \( \langle H(E) \rangle \) is the average histogram.

The thermal average of a thermodynamic quantity \( A \) can be evaluated by\(^2\)
\[\langle A \rangle_T = \frac{1}{Z} \sum_E g(E)A \exp(-E/k_BT),\]
where \( Z \) is the partition function defined by
\[
Z = \sum_E g(E) \exp(-E/k_BT)
\]
The canonical distribution at any \( T \) can be calculated simply by
\[
P(E, T) = \frac{1}{Z} g(E) \exp(-E/k_BT)
\]
In this work, we consider a energy range of interest\(^4\)\(^5\)\(^6\) \((E_{\text{min}}, E_{\text{max}})\). We divide this energy range to \( R \) subintervals, the minimum energy of each subinterval is \( E_{\text{min}} \) for \( i = 1, 2, ..., R \), and maximum of the subinterval \( i \) is \( E_{\text{max}} = E_{\text{min}} + 2\Delta E \), where \( \Delta E \) can be chosen large enough for a smooth boundary between two subintervals. The WL algorithm is used to calculate the relative DOS of each subinterval \( (E_{\text{min}}, E_{\text{max}}) \) with the modiﬁcation factor \( f_{\text{final}} = \exp(10^{-9}) \) and ﬂatness criterion \( x\% = 95\% \). We reject the suggested spin ﬂip and do not update \( g(E) \) and the energy histogram \( H(E) \) of the current energy level \( E \) if the spin-ﬂip trial would result in an energy outside the energy segment. The DOS of the whole range is obtained by joining the DOS of each subinterval \( (E_{\text{min}} + \Delta E, E_{\text{max}} - \Delta E) \).

III. RESULTS

We performed runs with different random quenched bond disorders at each given \( p \) for bond-conﬁguration averaging. The Edwards-Anderson order parameter, the magnetization and the susceptibility are obtained by standard MC simulations with the equilibrating time of \( t_e = 2 \times 10^6 \) MC steps/spin and the averaging time of \( t_a = 4 \times 10^6 \) MC steps/spin. The energy, speciﬁc heat and energy histograms are obtained by the WL technique in order to detect with precision the latent heat\(^4\)\(^2\)\(^3\).\(^4\).

In Fig. 2, we show the case of \( p = 0 \) where the energy and the magnetization versus \( T \) undergo a well-known ﬁrst-order transition\(^3\)\(^4\)\(^5\)\(^6\) with a discontinuity at the transition temperature \( T_c \). Figure 3 shows a double peak of the energy distribution at \( T_c \) which conﬁrms the ﬁrst-order behavior.

The latent heat \( \Delta E \) is deﬁned by the energy separation of the two peaks in the energy distribution. When the transition is of second order, the energy is continuous, namely \( \Delta E = 0 \). Using the WL technique which is very eﬃcient to detect ﬁrst-order transitions, we calculate the energy histogram with various small values of \( p \). We show in Fig. 4 the latent heat versus \( p \). As seen, \( \Delta E \) is not zero for \( p \leq 0.011 \). For \( p > 0.011 \), the phase transition is continuous. The disappearance of the latent heat for such a small amount of bond disorder is striking. This result is in agreement with the RG result on the 3D BEG model\(^2\), with the conjecture of Cardy and Jacobsen\(^2\) for the 3D case (see Fig. 1 of their paper), and with predictions of other earlier works\(^2\)\(^0\)\(^2\)\(^2\).\(^2\). To our knowledge, the present work is the ﬁrst veriﬁcation in 3D of the disappearance of the latent heat with a tiny amount of bond disorder. There is thus a tricritical point at \( p_c \simeq 0.011 \), beyond which the transition is of second order. This result is in agreement with MC simulations on the 3D BEG model\(^2\)\(^9\) and site-dilute Potts model\(^2\)\(^0\). An interesting question arises on the nature of the second-order transition coming from the ex ﬁrst-order transition. Malakis et al.\(^2\)\(^9\) have shown that the critical behavior belongs to the 3D random Ising model in violation of the universal-
For the case of pure ferromagnet, i.e. $p = 1$, the transition is a second-order one as expected for a 3D Ising ferromagnet.

Let us consider intermediate values of $p$ where the spin-glass phase is expected. The spin-glass phase is determined by the Edwards-Anderson freezing order parameter defined as

$$Q = \frac{1}{N} \frac{1}{t_\alpha} \sum_{i} \sum_{t=t_e}^{t_e+\tau_\alpha} S_i(t)$$

The averaging time $t_\alpha$ is taken as long as possible (at least several millions of MC steps per spin), after an equivalent equilibrating time $t_e$. In the calculation of $Q$, we follow each spin during the time $t_\alpha$ while calculating its time-averaged value. At the end, we take an average over all spins. So, $Q$ expresses the degree of spin freezing, independent of spin configuration.

The magnetization is defined as

$$M = \frac{1}{N} \frac{1}{t_\alpha} \sum_{i} \sum_{t=t_e}^{t_e+\tau_\alpha} S_i(t)$$

Here, $M$ is averaged over all spins before the time averaging. In a state with randomly frozen spins, the magnetization $M$ is zero so that it cannot distinguish a frozen phase from the paramagnetic phase.

We show now in Fig. 5 the energy and the specific heat versus $T$ for $p = 0.5$ where the bond disorder is strongest.

The magnetization $M$ versus $p$ at $T = 0.5$ and $T = 1.0$ is shown in Fig. 6. If a spin-glass phase is defined as the one where the magnetization is zero but not $Q$ at low $T$, then we find that the spin-glass phase exists in the range $0.055 < p < 0.62$.

The phase diagram $(T_c, p)$ is shown in Fig. 7 where $T_c$ is obtained at each $p$ when $Q$ goes down to zero. Note that the maxima of the specific heat and the susceptibility are located at a temperature higher than $T_c$. For example, at $p = 0.5$, $Q = 0$ at $T = T_c = 2 \pm 0.03$ while the peaks of $C$ and $\chi$ are at $T \approx 2.6$. There are several striking features shown in Fig. 7:

i) The phase diagram is not symmetric with respect to $p$. This is due to the fact that the FCC lattice is fully frustrated at $p = 0$ (antiferromagnetic) and non frustrated at $p = 1$. Note that the case where the pure system is non frustrated such as the simple cubic ferro- or antiferromagnets, the phase diagram $(T_c, p)$ is symmetric because the system is invariant with respect to the local transformation of every spin pair ($J_{ij} \rightarrow -J_{ij}, S_i \rightarrow -S_i$).
ii) The first-order transition line is terminated at $p = 0.011$ (see inset). As said earlier, in the present 3D model, such a tiny bond randomness suffices to change the first-order transition into a continuous transition.

iii) The long-range antiferromagnetic ordering is destroyed for $p \geq 0.055$ (we looked at $M$ down to $T \simeq 0$). The fact that such a very small $p$ suffices to destroy the long-range antiferromagnetic order is not a real surprise if we examine the nature of the long-range ordering in the antiferromagnetic limit: it is known that the ground-state (GS) spin configuration of the FCC antiferromagnet is a stacking of corner-sharing tetrahedra. Each tetrahedron has six choices of two spins up and two spins down on its corners (see Fig. 1). Such a stacking gives rise to an infinite number of GS configurations which are random with no long-range ordering except for 6 configurations composed of 3 configurations by stacking of up-spin planes and down-spin planes, alternately in $x$, $y$ or $z$ direction and 3 configurations due to the global reversal of all spins. When $T$ increases from zero, the system selects these 6 long-range GS configurations. This phenomenon is known as the selection of “order by disorder”\(^{50,51}\) which has been numerically verified in a large number of systems\(^{37,39,40,52}\). The selection of the long-range order is due to the entropy term which makes an extremely small difference with the other nearly-degenerate random configurations. It is not therefore surprising that the introduction of even a very small amount of ferromagnetic bonds breaks immediately the long-range order.

iv) It is interesting to note that we have observed two regions of reentrant phase at $0.02 < p < 0.055$ and $0.65 < p < 0.68$, bounded at the upper limit by the dotted lines. In these regions, at a given $p$ the magnetization is small at low $T$ and becomes zero when crossing the dotted line, below the transition line determined by the vanishing Edwards-Anderson order parameter $Q$. The dotted lines are determined using the temperatures at which the magnetizations vanish for different $p$ as shown in Fig. 8. Note that, depending on the system, there are several complex mechanisms leading to reentrant phases as seen in exactly solved Ising models\(^{53,54}\).

To study the relaxation of the freezing order parameter in the “steady” regime (after equilibrating time $t_e$), let us define the following quantities

$$Q_{\Delta t}(t) = \frac{1}{N} \frac{1}{\Delta t} \sum_i N \left| \sum_{t'=t}^{t+\Delta t} S_i(t') \right|,$$

$$\overline{Q_{\Delta t}} = \frac{1}{N t} \sum_t Q_{\Delta t}(t),$$

where $\sum_{t'=t}^{t+\Delta t}$ indicates a sum over $\Delta t$ MC steps per spin starting from $t$ and $\sum_{i}^{N}$ indicates a sum over the system.
sites. In Eq. (10), \( \overline{Q_{\Delta t}} \) is the average of \( Q_{\Delta t} \) calculated over all successive intervals of \( \Delta t \) steps until the end of the total run time \( t_a \). \( \overline{Q_{\Delta t}} \) is thus averaged with \( N_t = t_a/\Delta t = 10^4 \) intervals, taking \( t_a = 10^7 \) and \( \Delta t = 1000 \), for instance. Note that \( \Delta t \) is nothing but the so-called “waiting time” when measuring the system relaxation.

We emphasize that \( Q_{\Delta t}(t) \) is in fact the autocorrelation function. This is seen by writing the autocorrelation of the spin \( S_i \) after a waiting time \( \Delta t \) as

\[
|\langle S_i(t)S_i(t+\Delta t)\rangle| = \frac{1}{\Delta t} \left| \sum_{t'=t}^{t+\Delta t} S_i(t)S_i(t') \right|
\]

\[
= \frac{1}{\Delta t} \left( \pm \sum_{t'=t}^{t+\Delta t} S_i(t') \right) \quad \text{(since } S_i(t) = \pm 1) \]

\[
= \frac{1}{\Delta t} \sum_{t'=t}^{t+\Delta t} S_i(t')
\]

from which we obtain \( Q_{\Delta t}(t) \) by summing over all spins. When \( \Delta t = t_a \) we have \( Q_{\Delta t} \equiv Q \).

As shown in Fig. 9, the magnetization \( M \) is zero at any \( T \) for \( p = 0.5 \). Therefore, it is impossible to calculate the susceptibility defined by the variance of \( M \). However, we can do that for \( Q_{\Delta t} : \chi = N(\langle Q_{\Delta t}^2 \rangle - \langle Q_{\Delta t} \rangle^2)/T \). Of course, the result of \( \chi \) will depend on \( \Delta t \). We show an example of \( \chi \) for \( \Delta t = 30000 \) in Fig. 9. As will be seen below, this waiting time is much longer than the relaxation time at \( T \) higher than, but not too close to, \( T_c \simeq 2 \). Therefore, \( \chi \) shown in Fig. 9 is considered as time-independent in this temperature range with its peak at \( T \simeq 2.6 \).

We use the following SER function defined by

\[
Q_t = A \exp \left[ -(t/\tau)^b \right] + a,
\]

where \( t \) is the waiting time which is \( \Delta t \) in our definition given above, \( b \) the SER exponent, \( A \) a temperature-dependent constant, and \( \tau \) the SER relaxation time. Note that this definition, without the constant \( a \), has been used by most of previous authors\(^7\,13\,14\,17\,18\). They did not use the pre-factor \( t^{-b} \) as in Eq. (1). We have introduced \( a \) in order to use the infinite-time limit for the fitting because we have taken \( t \) from 1000 to 30,000 in the simulation which are rather short. At the infinite time limit, \( a \) is zero for \( T > T_c \), and \( a = Q \) for \( T < T_c \). We can assimilate \( Q \) calculated during several millions of MC steps to \( Q_{t=\infty} \). Note that doubling the averaging time \( t_a \) reduces the fluctuations of \( Q_{t=\infty} \) but does not significantly change its mean value at least up to the fifth digit, while the exponential term is much smaller, of the order of \( t^{-10} \). In our simulations, we take the equilibrating time \( t_e = 4 \times 10^6 \) MCs/spin and the averaging time \( t_a = 2 \times 10^6 \) MCs/spin. We plot in Fig. 10 the autocorrelation \( Q_{\Delta t}(t) \) for \( p = 0.5 \) as a function of \( t \) at \( T = 1.5 \), 1.7, ..., 3.3 obtained from the simulation. Fitting these curves with the formula (11) and using \( Q_{t=\infty} \), we get the parameters \( A \), \( a \) and \( b \) and \( \tau \) which are shown in Tab. I. Note that, the SER function in Ref. 14 and 18 was studied only for \( T > T_c \).

The dependence of \( a \), \( b \), \( A \) and \( \tau \) on \( T \) is shown in Figs. 11 and 12. Several remarks are in order:
FIG. 10. \( Q_{\Delta t}(t) \) vs \( t \) (in unit of MC step) at \( T = 1.5, 1.7, \ldots, 3.3 \) (from top to bottom), for \( p = 0.5 \).

![Graph](image)

FIG. 11. Parameters \( a \) (top) and \( b \) (middle), and \( A \) (bottom) as functions of \( T \). See text for comments.

![Graph](image)

FIG. 12. SER relaxation time \( \tau \) as a function of \( T \).

![Graph](image)

TABLE I. Values of \( a, b, A, \) and \( \tau \) at several \( T \)

<table>
<thead>
<tr>
<th>( T )</th>
<th>( \tau )</th>
<th>( b )</th>
<th>( a )</th>
<th>( A )</th>
</tr>
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<tbody>
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<td>0.246619</td>
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i) \( a \) decreases with increasing \( T \). However it goes to zero only far from \( T_c \), namely deep inside the paramagnetic phase. The residue value after \( T_c \) is due to the finite size effect and to short-range correlations which exist at \( T > T_c \).

ii) \( A \) changes the curvature at \( T_c \approx 2 \). If this coefficient is related to the pre-factor \( t^{-x} \) of Ogelski\(^6\), then the change of curvature of \( x \) at the glass transition temperature that he observed may be due to the same physical origin.

iii) \( b(T) \) changes the curvature at \( T_c \approx 2 \) where it is equal to \( \approx 0.33 \). Note that the determination of \( b \) at a given \( T \) does not depend on \( T_c \), unlike the critical exponents. The value \( b = 1/3 \) at \( T_c \) is what predicted by theories\(^7,10,13\). It is also what was found numerically for different models of Ising SG: \( \pm J \) on simple cubic lattice\(^6\) and Sherrington-Kirkpatrick model\(^14\). This value is also close to what experimentally found in various types of SG, not necessarily of Ising models\(^15-18\). We believe that at least at \( T_c, b = 1/3 \) is universal. The value of \( b \) increases with increasing \( T \). Note however that we did not go very far into the paramagnetic phase (\( T_c < T < 3.3 \)) so that \( b \) is far from the expected value \( b = 1 \) for the true paramagnetic phase (\( T \gg T_c \)), but the tendency towards \( 1 \) is clearly seen in the curve of \( b \) of Fig. 11.

iv) The SER relaxation time \( \tau \) becomes extremely large at \( T \approx 2 \) as expected from the critical slowing down.

To close this section, let us show the size dependence for some quantities. Figure 13 shows the specific heat and the susceptibility versus \( T \) with system sizes \( N = 4 \times L^3 \) where \( L = 12, 16 \) and 20, for \( p = 0.5 \).

We see that the size effect is not strong at the strongest disorder \( p = 0.5 \). Our results for the SER shown above with \( p = 0.5 \) and \( L = 12 \) will not be altered with larger
IV. CONCLUDING REMARKS

In this paper, we obtained three striking results:

(i) the disappearance of the latent heat when a very small amount of bond disorder (\( \simeq 1\% \)) is introduced into our system. This result is in agreement with theories\(^{21-23,28}\) and simulations on 3D models\(^{29,30}\) having a first-order transition in their pure state.

(ii) the long-range ordering of the pure system considered here (FCC antiferromagnet) is lost at a very small amount of bond disorder (\( \simeq 5\% \)). This means that the entropy created by the so-called “order by disorder” which is at the origin of the long-range ordering in the FCC antiferromagnet is so small that it is overwhelmed by just a small bond disorder.

(iii) the SER is verified with \( b = 1/3 \) at the freezing transition in agreement with MC simulations on other Ising SG\(^{6,14}\) and with theories\(^{7,13}\). This value is close to the ones experimentally observed in various SG\(^{15-18}\). We believe that the value \( b = 1/3 \) at \( T^* \) in the SER is universal for different kinds of SG.

Finally, we note that we studied here a dozen bond configurations of size \( 24 \times 21^3 \). At this large number of bonds, statistical fluctuations between samples are almost zero. We did not study therefore the size effect on the relaxation time in the present paper, reserving a priority for simulations using the whole range of ferromagnetic bond concentration which took an enormous CPU time. In addition, we did not study SG exponents such as \( \nu \) and \( z \) (dynamic exponent) so the finite-size scaling is not necessary at this stage. We believe however that our results on the phase diagram and the relaxation time remain valid for larger sizes. Size effects in SG are certainly much smaller than in pure systems because of the existence of different correlation lengths at the glass transition due to different degrees of local disorders, in the same manner as the existence of many different exponential relaxation times whose superposition leads to the slow SER. It is however very interesting to study the finite-size scaling to calculate the universality class of a system having a first-order transition after its crossover to a second-order due to the introduction of a random bond disorder. This problem is left for a future study.

ACKNOWLEDGMENTS

The authors are grateful to A. Nihat Berker and H. Orland for helpful comments and suggestions.

V. Thanh Ngo would like to thank the University of Cergy-Pontoise for a financial support during the course of this work. He is grateful to Vietnam National Foundation for Science and Technology Development (Nafosted) for support (Grant No. 103.02-2011.55).


