Fluctuation effects in the Ising model with reduced interaction and quenched disorder

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The effect of quenched random fields and local perturbations of critical temperature on the critical behavior at phase transitions is studied within the framework of an exactly solvable model that takes into account interaction of fluctuations with equal and opposite momenta. Using the replica method the dimensional reduction by 2 for systems with finite-range interaction and quenched random fields is explicitly shown. For interaction of the infinite range the model demonstrates the mean-field critical asymptotics independently of dimensionality or the presence of random fields.

I. THE MODEL

Within the past two decades there has been considerable interest in critical behavior of random systems.[1–9] The renormalization-group (RG) treatment does not provide an unambiguous solution of this problem.[1–7] Therefore, it is very instructive to consider phase transitions in random systems within the framework of exactly solvable models that take into account fluctuation effects at least partly. In Refs. 8–11 the spherical model[12] has been applied for the consideration of the instability of an ordered phase due to a random field. In the present paper we study effects of quenched random fields and local temperature fluctuations in the reduced ϕ^4 model which takes into account interaction of fluctuations with equal and opposite momenta only. This model allows for an exact solution. It is very convenient for investigation of critical phenomena in complex systems. Even though in its simplest version it has the same critical asymptotics as the spherical model,[13,14] it explicitly demonstrates major qualitative results that have been obtained within renormalization-group theory, including fluctuation induced first-order phase transitions.[15,16] Renormalization-group theory and the model lead to a similar critical behavior for phase transitions in orthorhombic high-T_c superconductors with d pairing[17] and for oxygen ordering near a structural phase transition in Y-Ba-Cu-O.[18] Below we briefly review basic features of the "pure" model and derive relationships that will be used in subsequent sections for consideration of the disordered case.

We start from the Ginzburg-Landau functional with a scalar order parameter $\phi(\mathbf{x})$,

$$H=\tfrac{1}{2}\int d^dx[\tau\phi^2(\mathbf{x})+c(\nabla\phi(\mathbf{x}))^2+\tfrac{1}{4}g\phi^4(\mathbf{x})-2h\phi(\mathbf{x})], \tag{1}$$

where $\tau\propto T-T_c$, T_c is a trial critical temperature and h is a constant field. To calculate the partition function with the functional (1) exactly, we split interaction terms as follows:

$$\int d^dx\,\phi^4(\mathbf{x})\rightarrow\frac{1}{V}a^2[\phi],\quad a[\phi]=\int d^dx\,\phi^2(\mathbf{x}), \tag{2}$$

where V is the system volume. After such a reduction the model takes into account interaction of fluctuations with equal and antiparallel momenta only. This can be seen if one rewrites Eq. (2) in the momentum representation. Then the reduction (2) is equivalent to splitting the δ function, which provides momentum conservation, into the product of two δ functions: $\delta(\mathbf{q}_1+\mathbf{q}_2+\mathbf{q}_3+\mathbf{q}_4)\rightarrow\delta(\mathbf{q}_1+\mathbf{q}_2)\delta(\mathbf{q}_3+\mathbf{q}_4)$. This reduction transfers the ϕ^4 model into the universality class of the spherical model.[13,14]

In order to calculate a functional integral with respect to $\phi(\mathbf{x})$ one has to transform the fourth-order form with respect to ϕ in the Boltzmann factor in the partition function into a bilinear form. This can be done with the help of a transformation analogous to that of Hubbard-Stratonovich,

$$\exp\left[-\frac{V}{2}K\left(\frac{a[\phi]}{V}\right)\right]=\frac{1}{2\pi}\int dx\,dy\,\exp\left[-\frac{V}{2}K\left(\frac{x}{V}\right)+i(xy-ya[\phi])\right], \tag{3}$$

where K is an arbitrary function. After applying the transformation (3) Gaussian integrals with respect to ϕ in the partition function can be calculated. As a result, the partition function takes the form:

$$Z\propto\int_{-\infty}^{\infty}dx\,dy\,\exp\left\{-\frac{V}{2}\left[\tau x+\frac{1}{4}gx^2-xy+\frac{1}{V}\sum_q\ln|c\mathbf{q}^2+y|-\frac{h^2}{y}\right]\right\}. \tag{4}$$

The summation of the kind $\Sigma_{\mathbf{q}}\ln|c\mathbf{q}^2+y|$ has to be cutoff since it diverges on the upper limit. However, critical asymptotics should not depend upon momentum cutoff. For the spacial dimension $2<d<4$ this can be handled by renormalizing the summation and then setting the momentum cutoff to be equal to infinity.[14,15] For $d\geqslant4$ the sum becomes nonrenormalizable and we must maintain the momentum cutoff explicitly. As we demonstrate below the dependence upon the momentum cutoff is absorbed into a renormalization of the trial value of the critical temperature and into an insignificant constant addition to the free energy. If the cutoff momentum is equal to Λ, then

$$\sum_{\mathbf{q}}\ln|c\mathbf{q}^2+y|=[y\theta_d(\Lambda,c)-f_d(y;c)]V, \tag{5}$$

where

$$\theta_d(\Lambda;c)=\frac{S_d}{(2\pi)^d}\times\begin{cases}\Lambda^{d-2}/c(d-2), & d\neq 2\\ [1+\ln(c\Lambda^2)]/2c, & d=2,\end{cases}$$

$$f_d(y;c)=\frac{S_d}{(2\pi)^d}\times\begin{cases}\dfrac{\pi y^{d/2}}{dc^{d/2}\sin[\pi(d-2)/2]}\equiv\kappa(c)y^{d/2}, & \\ & d\neq \text{even}\\ -\dfrac{1}{d}\left(-\dfrac{y}{c}\right)^{d/2}\ln|y|\equiv-\mu(c)y^{d/2}\ln|y|, & \\ & d=\text{even},\end{cases}\tag{6}$$

and S_d is the surface area of a d-dimensional unit radius sphere. $\theta_d(\Lambda)$ is used to renormalize $x,x\rightarrow x+\theta_d(\Lambda)$, which consequently results in the renormalization of the trial critical temperature $\tau,t=\tau+g/2\theta_d(\Lambda)$. As a result the partition function becomes

$$Z\propto\int dx\;dy\;\exp[-VF(x,y)];$$

$$F(x,y)=\tfrac{1}{2}[x(t-y)+gx^2/4-f_d(y,c)-h^2/y].\tag{7}$$

In the thermodynamic limit, $V\rightarrow\infty$, one can exactly calculate integrals in Eq. (7) using the method of the steepest descend. Hence, all thermodynamic quantities can be calculated. The free-energy density is given by Eq. (7) with x and y being a solution of the system of equations: $\partial F/\partial x=0$, $\partial F/\partial y=0$. After eliminating x from these equations we derive an equation for y,

$$t-y+(g/2)[(h/y)^2-f_d'(y;c)]=0.\tag{8}$$

An equilibrium value of the order parameter is given by

$$\phi_0=-\partial F/\partial h=h/y(h)\tag{9}$$

with y determined by Eq. (8). When $h=0$ a nontrivial real value of the order parameter, $\phi_0\neq 0$, exists for any $d>2$ and $t<0$. Equations (8) and (9) give $\phi_0=(-2t/g)^{1/2}$. When $d\leqslant 2$ a nontrivial solution, $\phi_0\neq 0$, does not exist. Therefore, $d=2$ is a lower marginal dimension for the model. The critical exponent β is equal to $\frac{1}{2}$. The model gives more interesting results when calculating the critical exponent δ. When $t=0$ and $h\neq 0$ Eqs. (8) and (9) reduce to the equation for the order parameter

$$-\frac{h}{\phi_0}+\frac{g}{2}\phi_0^2-\frac{g\kappa(c)d}{4}h^{(d-2)/2}\phi_0^{(2-d)/2}=0,\quad d\neq\text{even}\tag{10a}$$

$$-\frac{h}{\phi_0}+\frac{g}{2}\phi_0^2+\frac{g\mu(c)}{2}\left(\frac{h}{\phi_0}\right)^{(d-2)/2}\left[1+\frac{d}{2}\ln\left(\frac{h}{\phi_0}\right)\right]=0,$$

$$d=\text{even}.\tag{10b}$$

In the limit $h\rightarrow 0$ the solutions of Eqs. (10) are easily found:

$$\phi_0\approx\begin{cases}(\kappa(c)d/2)^{2/(d+2)}h^{(d-2)/(d+2)}, & 2<d<4\\ [-(h/3c^2)\ln(hc)]^{1/3}, & d=4\\ (2/g)^{1/3}h^{1/3}, & d>4.\end{cases}\tag{11}$$

From Eq. (11) it follows that $d=4$ is an upper marginal dimension: when $2<d<4$ the critical exponent δ coincides with the RG result, $\delta=(d+2-\eta)/(d-2+\eta)$, if numerically small exponent η is set to be equal to zero, for $d>4$ the model gives the mean-field result, $\delta=3$, and for $d=4$ a logarithmic correction to δ arises.

Furthermore, for $2<d<4$ a crossover effect from critical behavior to mean field when h increases can be obtained from Eq. (10a). In other words, when $2<d<4$ and

$$h\gg g^{(d+2)/2(4-d)}\kappa(c)^{3/(4-d)}\tag{12}$$

the first term in Eq. (10a) becomes dominant and the critical exponent δ changes from $(d+2)/(d-2)$ to 3. Since c is a squared correlation radius far away from the critical point $(t\sim 1)$, the Ginzburg parameter Gi, which defines the width of the fluctuation region, is proportional to $Gi\propto c^{-d/(4-d)}\propto\kappa(c)^{2/(4-d)}$. Therefore, it follows from Eq. (12) that the crossover from critical to mean-field behavior occurs when $h\gg Gi^{3/2}$. This coincides with the results of the RG analysis.

II. RANDOM FIELD

In this section we examine an effect of a random field on phase transitions within the context of the exactly solvable model. The Ginzburg-Landau functional with a scalar order parameter $\phi(\mathbf{x})$ has the form

$$H=\tfrac{1}{2}\int d^dx[\tau\phi^2(\mathbf{x})+c(\nabla\phi(\mathbf{x}))^2+\tfrac{1}{4}g\phi^4(\mathbf{x})-2h\phi(\mathbf{x})$$
$$-2h(\mathbf{x})\phi(\mathbf{x})],\tag{13}$$

where h and $h(\mathbf{x})$ are constant and random fields, respectively. To obtain the free energy averaged over the random field, following Ref. 19, we replicate n times the partition function Z by defining an n-component vector $\boldsymbol{\phi}(\mathbf{x})\equiv[\phi_1(\mathbf{x}),...,\phi_n(\mathbf{x})]$. Hence,

$$F=-\frac{\partial}{\partial n}\left[\int D^n\phi(\mathbf{x})\exp(-H_{\text{eff}}[\boldsymbol{\phi}(\mathbf{x})])\right]_{n=0},\tag{14}$$

where $H_{\text{eff}}[\boldsymbol{\phi}]$ is given by

$$H_{\text{eff}}[\phi(\mathbf{x})]\equiv\frac{1}{2}\int_{-\infty}^{\infty}d^dx\Bigg\{\tau|\phi(\mathbf{x})|^2+c(\nabla\phi(\mathbf{x}))^2-Q[\boldsymbol{\phi}]$$
$$+\sum_{i=1}^{n}\left[\frac{1}{4}g\phi_i^4(\mathbf{x})-2h\phi_i(\mathbf{x})\right]\Bigg\}\tag{15}$$

with

$$Q[\boldsymbol{\phi}]=\int d^dx\,\ln\Bigg\{\frac{1}{\sqrt{2\pi B}}\int dh(\mathbf{x})\exp\left[-\frac{h^2(\mathbf{x})}{2B}\right]$$
$$\times\exp\left[\sum_{i=1}^{n}h(\mathbf{x})\phi_i(\mathbf{x})\right]\Bigg\},\tag{16}$$

where we supposed that $h(\mathbf{x})$ is a δ-correlated Gaussian random variable, $\langle h(\mathbf{x})\mathrm{h}(\mathbf{x}')\rangle=B\delta(\mathbf{x}-\mathbf{x}')$. Then $H_{\text{eff}}[\boldsymbol{\phi}]$ becomes

$$H_{\text{eff}}=\frac{1}{2}\int d^dx\Bigg[\tau|\phi(\mathbf{x})|^2+c(\nabla\boldsymbol{\phi}(\mathbf{x}))^2+\sum_{i=1}^{n}\left[\frac{1}{4}g\phi_i^4(\mathbf{x})\right.$$
$$\left.-2h\phi_i(\mathbf{x})\right]-B\left(\sum_{i=1}^{n}\phi_i(\mathbf{x})\right)^2\Bigg].\tag{17}$$

Now, for the functional (17) we apply a reduction similar to Eq. (2),

$$\int d^d x \varphi_i^4(\mathbf{x}) \to \frac{1}{V} a^2[\varphi_i], \quad a[\varphi_i] = \int d^d x \varphi_i^2(\mathbf{x}), \tag{18}$$

so that the partition function takes the form

$$Z \propto \int \left[\prod_{i=1}^{n} D\phi_{i\mathbf{q}} dx_i dy_i \right] \exp\left\{ -\frac{V}{2} \sum_{i=1}^{n} \left[\tau x_i + \frac{1}{4} g x_i^2 - x_i y_i - \frac{2}{\sqrt{V}} h \phi_{i0} \right] - \frac{1}{2} \sum_{i=1}^{n} \sum_{\mathbf{q}} |\phi_{i\mathbf{q}}|^2 (y_i + c\mathbf{q}^2 - B) + \frac{B}{2} \sum_{\mathbf{q}} \sum_{\substack{i,j=1 \\ i \neq j}}^{n} \phi_{i\mathbf{q}} \phi_{j-\mathbf{q}} \right\}. \tag{19}$$

In order to calculate functional integrals in Eq. (19) the form in the exponent must be diagonalized with respect to components of the vector $\boldsymbol{\phi}$. After diagonalization we require that $y_i = y_j \equiv y$ since only this choice reproduces the pure ϕ^4 upon suppression of the random field. Explicitly, this is so because in the limit of $B \to 0$ the degeneracy of the eigenvalues of every other choice does not reduce to n fold as expected from considerations of the simple ϕ^4 model treated within the context of the replica method. In this case the $n \times n$ matrix of interest has only two distinct eigenvalues,

$$\lambda_1 = -\tfrac{1}{2}(y + c\mathbf{q}^2); \quad \lambda_2 = -\tfrac{1}{2}(y + c\mathbf{q}^2 - nB), \tag{20}$$

where λ_1 is $(n-1)$-fold degenerate. These eigenvalues can be used to diagonalize the $n \times n$ matrix with respect to ϕ_i. As the result one can calculate functional integrals in Eq. (19) and arrive at

$$Z \propto \int_{-\infty}^{\infty} \left(\prod_{i=1}^{n} dx_i \right) dy \, \exp\left\{ -\frac{V}{2} \sum_{i=1}^{n} \left[\tau x_i + \frac{1}{4} g x_i^2 - y x_i \right] + \frac{(1-n)}{2} \sum_{\mathbf{q}} \ln|y + c\mathbf{q}^2| - \frac{1}{2} \sum_{\mathbf{q}} \ln|y + c\mathbf{q}^2 - nB| + \frac{nVh^2}{2(y-nB)} \right\}. \tag{21}$$

After treating the summations with respect to $\mathbf{q}$ as in the previous section we derive

$$2F(x_i, y, h) = \sum_{i=1}^{n} \left[t x_i + \frac{1}{4} g x_i^2 - y x_i \right] + (1-n) f_d(y;c) - f_d(y - nB; c) - \frac{nh^2}{y - nB}. \tag{22}$$

Using $\partial F / \partial x_i = 0$, we simplify the above equation by replacing x_i

$$2F(y,h) = \sum_{i=1}^{n} \left[-\frac{t^2}{g} - \frac{y^2}{g} + \frac{2yt}{g} - f_d(y,c) - \frac{h^2}{y - nB} \right] + f_d(y;c) - f_d(y - nB; c) \tag{23}$$

and the function $y(h)$ can be found from the equation

$$\partial F(y,h)/\partial y = 0 \Rightarrow -2ny/g + 2nt/g - (n-1) f_d'(y;c) - f_d'(y - nB; c) + nh^2/(y - nB)^2 = 0. \tag{24}$$

The solution for $y = y(h)$ must then be substituted in $F(y,h)$ and the disorder averaged value of the equilibrium free energy is given by

$$F(h) = \lim_{n \to 0} [F(y(h), h)/n]. \tag{25}$$

An expression of an averaged order parameter, ϕ_0, is given by

$$\phi_0 = \lim_{n \to 0} (1/n) [-\partial F(y,h)/\partial h]_{y = y(h)} = \lim_{n \to 0} [h/y(h)]. \tag{26}$$

To find ϕ_0 from Eqs. (24) and (26) one has to expand Eq. (24) in powers of n up to the lowest order of n. For d not even (including nonintegers) and d even the resulting equations for the order parameter are

$$-\frac{2}{g}\left(\frac{h}{\phi_0}\right) + \frac{2t}{g} + \phi_0^2 - \frac{\kappa(c)}{2}\left(\frac{h}{\phi_0}\right)^{(d-2)/2} + \frac{\kappa(c)(d-2)B}{4}\left(\frac{h}{\phi_0}\right)^{(d-4)/2} = 0, \quad d = \text{not even} \tag{27a}$$

$$-\frac{2}{g}\left(\frac{h}{\phi_0}\right) + \frac{2t}{g} + \phi_0^2 + \frac{\mu(c)d}{2}\left(\frac{h}{\phi_0}\right)^{(d-2)/2} \ln\left(\frac{h}{\phi_0}\right) + \mu(c) \times \left(\frac{h}{\phi_0}\right)^{(d-2)/2} - \mu(c)B(d-2)\left(\frac{h}{\phi_0}\right)^{(d-4)/2} \left[\frac{1}{2} + \frac{d}{2(d-2)} + \frac{d}{4}\ln\left(\frac{h}{\phi_0}\right)\right] = 0, \quad d = \text{even}. \tag{27b}$$

In the limit of $B \to 0$ Eqs. (23)–(27) reproduce the corresponding ones of the "pure" ϕ^4 model.

From Eqs. (27) it is seen that for $d \leqslant 4$ no solution for ϕ_0 exists. Thus, the random field, regardless of how weak it is, destroys a long-range order for $d \leqslant 4$. When $d > 4$ we obtain the second order transition occurring at $t = 0$ and $h = 0$. When $4 < d < 6$ the third and the fourth terms in Eq. (27a) can be omitted because they become of smaller order in h compared to the last term. Therefore, this equation gives the same critical asymptotics as the "pure" model with the lower dimension $2 < d' (= d-2) < 4$. Namely, for critical exponents β and δ we have $\beta = \frac{1}{2}$, $\delta = (d'+2)/(d'-2)$. When $d > 6$ in the limit $h \to 0$ only the first three terms in Eqs. (27) survive,

$$-(2/g)(h/\phi_0) + 2t/g + \phi_0^2 = 0, \tag{28}$$

and we arrive at the mean-field critical exponents. The critical exponent δ for the random d-dimensional system is exactly the same as that of a $(d-2)$-dimensional pure system. In addition, as it follows from Eq. (27b), the random six-dimensional system, has the same logarithmic corrections as the pure four-dimensional model. So, the model explicitly demonstrates the same dimensional crossover in the presence of quenched random fields found in the RG approach. It may seem contradictory that the model with the one-component order parameter has a lower critical dimension $d_c=4$. Indeed, the functional (13) corresponds to the random-field Ising model which has $d_c=2$. However, after the reduction (18) the model belongs to the spherical model universality class and, therefore, has symmetry $O(N=\infty)$.[20] Phase transitions in the spherical model with random fields have been considered in Ref. 8 where the condensation of the ideal Bose gas in the presence of random sources has been studied. The ideal Bose-Einstein condensation at constant volume is identical to a phase transition in the spherical model. Results obtained in Ref. 8 are analogous to the results of our model.

The crossover from critical to the mean-field behavior in $4<d<6$ and nonzero random field occurs when

$$h \gg g[B\kappa(c)]^{3/(6-d)}. \tag{29}$$

The above inequality is always true when the fluctuations are suppressed in the limit of $c\to\infty(\kappa\to 0)$. Moreover, when $\kappa(c)=0$ and $\mu(c)=0$ Eqs. (27) reduce to Eq. (28). That means that in systems with long-range interactions a second-order phase transition is restored and critical exponents become the same as the mean-field ones for any dimensionality regardless of the presence of random fields.

III. RANDOM TEMPERATURE

We now examine the problem of frozen in nonmagnetic impurities that cause a random perturbation of local temperature. Phase transitions in such a system can be described by a free-energy functional

$$H=\frac{1}{2}\int d^dx[\tau_0\phi^2(\mathbf{x})+\tau(\mathbf{x})\phi^2(\mathbf{x})+c(\nabla\phi(\mathbf{x}))^2+\tfrac{1}{4}\,g\phi^4(\mathbf{x}) -2h\phi(\mathbf{x})], \tag{30}$$

where $\tau(\mathbf{x})$ is a δ-correlated random function, $\langle\tau(\mathbf{x})\tau(\mathbf{x}')\rangle=B\delta(\mathbf{x}-\mathbf{x}')$, with zero mean value. After following the steps of the replica method and treating the problem within the context of the exactly solvable model we derive,

$$2F(x_i,y_i,h)=\sum_{i=1}^{n}\left[tx_i+\frac{1}{4}gx_i^2-x_iy_i-f_d(y_i;c)-\frac{h^2}{y_i}\right] -B\left(\sum_{i=1}^{n}x_i\right)^2. \tag{31}$$

Using the saddle-point equations $\partial F/\partial x_i=0$ and $\partial F/\partial y_i=0$, we obtain y_i and x_i which are then used in Eqs. (25) and (31) to find the averaged free energy. Having in mind that the only physical choice is $y_i=y_j\equiv y$, which implies $x_i=x_j\equiv x$ the saddle-point equations give

$$t-y+\frac{g}{2}\left[\frac{h^2}{y^2}-f'_d(y,c)\right]+2nBf'_d(y;c)-\frac{2nBh^2}{y}=0. \tag{32}$$

In the limit $n\to 0$ Eq. (32) is equivalent to Eq. (8). So, critical asymptotics in this case are the same as those given by the model in the pure case. This is in agreement with the Harris criterion[21] according to which disorder of the random temperature type changes the critical behavior only if the critical exponent α of the pure system is positive. The critical exponent α in our model is $\alpha=-(4-d)/(d-2)$ that is negative for $d<4$. For $d>4$ the model gives the mean-field behavior which is also insensitive to a randomness.

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