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A unified picture of ferromagnetism, quasi-long range order and criticality in random field models

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By applying the recently developed nonperturbative functional renormalization group (FRG) approach, we study the interplay between ferromagnetism, quasi-long range order (QLRO) and criticality in the \( d \)-dimensional random field \( O(N) \) model in the whole \((N, d)\) diagram. Even though the "dimensional reduction" property breaks down below some critical line, the topology of the phase diagram is found similar to that of the pure \( O(N) \) model, with however no equivalent of the Kosterlitz-Thouless transition. In addition, we obtain that QLRO, namely a topologically ordered "Bragg glass" phase, is absent in the 3-dimensional random field XY model. The nonperturbative results are supplemented by a perturbative FRG analysis to two loops around \( d = 4 \).

How the phase behavior and ordering transitions of a system are affected by the presence of a weak random field remains in part an unsettled problem. Heuristic and rigorous arguments show that the lower critical dimension below which no long-range order is possible is 2 for the random field Ising model (RFIM) [1, 2] and 4 for models with a continuous symmetry (RFO(N)M with \( N > 1 \)) [3, 4]. However, this leaves aside two questions: first, the nature of the critical behavior in random field models, a question connected to the breakdown of the so-called "dimensional reduction" (DR) property that relates the critical exponents of the RFO(N)M to those of the pure \( O(N) \) model in two dimensions less [2]; and second, the possible occurrence of a low-temperature phase with quasi-long range order (QLRO), i.e., a phase characterized by no magnetization and a power-law decrease of the correlation fuctions, in models with a continuous symmetry [5, 6]. Progress has been made to better circumcise this latter point. It has indeed been shown that QLRO is absent for \( N = 2 \) when disorder is strong and for \( N > 3 \) for arbitrarily weak random field [7]; but this still keeps open the cases \( N = 2, 3 \) in the physical dimensions \( d = 2, 3 \).

Those questions are important because, on top of purely theoretical motivations, they concern the behavior of the known experimental realizations of random field models. This is the case for instance of vortex lattices in disordered type-II superconductors [8, 9]. In such systems, the randomly pinned lattice of vortices can be mapped onto an "elastic glass model" [7, 9], whose simulated structure is similar to that of the pure \( O(N) \) model, with however no equivalent of the Kosterlitz-Thouless (KT) transition.

Our starting point is the standard effective hamiltonian for the RFO(N)M in \( d \) dimensions with an \( N \)-component field \( \chi(x) \) and uncorrelated random fields taken from a Gaussian distribution with zero mean and variance \( \Delta \). After introducing \( n \) replicas in order to perform the average over quenched disorder (taking at the end the limit \( n \to 0 \)), it can be rewritten as:

\[
S\{\{\chi_a\}\} = \int_x \left\{ \frac{1}{2T} \sum_{a=1}^n \left[ \partial_x \chi_a \cdot \partial_x \chi_a + \tau \chi_a^2 + \frac{n}{12} \chi_a^4 \right] - \frac{\Delta}{2T^2} \sum_{a,b=1}^n \chi_a \cdot \chi_b - \sum_{a=1}^n J_a \cdot \chi_a \right\},
\]

where we have introduced sources acting on each replica separately, which therefore explicitly break the permutation symmetry between the replicas.

We apply in this work the nonperturbative FRG formalism for the RFO(N)M recently proposed by us [12]. To keep the presentation short but sufficiently self-contained, we first sketch the main steps of the approach. It is based on an exact RG equation for the effective average action \( \Gamma_k\{\{\phi_a\}\} \). \( \Gamma_k \) interpolates between the bare action, eq. (1), and the usual effective action (i.e., the Legendre transform of the partition function associated with eq. (1)) as the running scale \( k \) moves from microscopic \((k = \Lambda)\) to macroscopic \((k \to 0)\) scale. It is built by integrating out fluctuations with momenta larger than
$k$ and its flow obeys an exact equation,
\[
\partial_t \Gamma_k[\{\phi_a\}] = \frac{1}{2} \text{Tr}(\partial_t R_k(q)(\Gamma^{(2)}_k[\{\phi_a\},q] + \mathbb{I} R_k(q))^{-1})
\] (2)
where $\partial_t$ is a derivative with respect to $t = \ln(k/\Lambda)$, $\Gamma^{(2)}_k$ is the tensor formed by the second functional derivatives of $\Gamma_k$ with respect to the fields $\phi_a(q)$. $\mathbb{I}$ is the unit tensor, and the trace involves an integration over momenta as well as a sum over replica indices and N-vector components. $R_k(q)$ is the infrared cutoff introduced to suppress the low-momentum modes. In practice, solving eq. (2) numerically requires the introduction of approximation schemes which amount to truncate the functional form of $\Gamma_k$. Guided by the physics of the problem at hand and by the mounting work on the method [13], one can then formulate a nonperturbative RG description. For the RFO(N)M we have argued that the two main ingredients allowing to study the long-distance physics are (i) the derivative expansion, which approximates the momentum dependence of the (1PI) vertex functions [13], and (ii) the expansion in increasing number of free replica sums, which is equivalent to include increasing-order cumulants of the renormalized distribution of the quenched disorder [12].

The simplest nonperturbative FRG description of the RFO(N)M relies on the following truncation:
\[
\Gamma_k[\{\phi_a\}] = \int_x \left\{ \sum_{a=1}^n \left( \frac{1}{2} Z_{m,k} \partial \phi_a \cdot \partial \phi_a + U_k(\phi_a) \right) - \frac{1}{2} \sum_{a,b=1}^n V_k(\phi_a, \phi_b) \right\}
\] (3)
with one single wavefunction renormalization for all fields, $Z_{w,k}$, evaluated at the field configurations that minimize the 1-replica potential $U_k$ (pseudo first-order derivative expansion [13]); the 1-replica $U_k$ and 2-replica $V_k$ parts of the effective potential are obtained from $\Gamma_k$ taken for uniform fields. Inserting eq. (3) into eq. (2), leads to coupled partial differential equations for the functions $U_k(\tilde{\rho})$ and $V_k(\tilde{\rho}, \tilde{\rho}', z)$, where $\tilde{\rho} = \frac{1}{2} \phi^2$ and $z = \phi \cdot \phi' / \sqrt{4 \tilde{\rho} \tilde{\rho}'}$, and to a running anomalous dimension defined as $\eta_k = -\partial_t \log Z_{m,k}$. The sought FPs being at zero temperature, one also introduces a running temperature $T_k$ and the associated exponent $\theta_k = \partial_t \log T_k$. This is most conveniently done by defining a renormalized disorder strength $\Delta_{m,k} = (2 \hat{\rho}_{m,k})^{-1} \partial_z V_k(\tilde{\rho}_{m,k}, \tilde{\rho}_{m,k}, z = 1)$, where $\hat{\rho}_{m,k}$ corresponds to the minimum of $U_k(\tilde{\rho})$. Defining then $\tilde{T}_k$ as $Z_{m,k} k^2 \Delta_{m,k}$, one can see that it reduces to the bare temperature $T$ at the microscopic scale $\Lambda$.

The flow equations can be expressed in a scaled form by introducing renormalized dimensionless quantities, $u_k(\tilde{\rho}) = T_k^{-d-4} U_k(\tilde{\rho})$, $v_k(\tilde{\rho}, \tilde{\rho}', z) = T_k^{-d-4} V_k(\tilde{\rho}, \tilde{\rho}', z)$, $\tilde{\rho} = Z_{w,k} T_k k^{-(d-2)} \tilde{\rho}$. An expression for $\eta_k$ is derived by considering the flow of the transverse component of $\Gamma^{(2)}_k$ evaluated for uniform fields and zero momentum:
\[
\eta_k = 8 C^-_d / [d \{ u_{pp}(\tilde{\rho}) - 4 \rho_m v_{pp'} | m_{d,1}^d(2 \rho_m u_{pp}, 0) + 2 \rho_m^2 v_{z} | m_{d,2}^d(2 \rho_m u_{pp}, 0) + m_{d,2}^d(2 \rho_m u_{pp}, 0) \}]
\] (4)
where $C^-_d$ is most conveniently done by defining a renormalized disorder strength $\Delta_{m,k} = (2 \hat{\rho}_{m,k})^{-1} \partial_z V_k(\tilde{\rho}_{m,k}, \tilde{\rho}_{m,k}, z = 1)$, where $\hat{\rho}_{m,k}$ corresponds to the minimum of $U_k(\tilde{\rho})$. Defining then $\tilde{T}_k$ as $Z_{m,k} k^2 \Delta_{m,k}$, one can see that it reduces to the bare temperature $T$ at the microscopic scale $\Lambda$.

The central results of our study are summarized on the phase diagram displayed in Fig. 1b. In region III there are no phase transitions, the RFO(N)M always stays disordered. In region I there is a transition to a ferromagnetic phase at a critical point governed by a zero-temperature nonanalytic FP. The nonanalyticity is strong enough (a “cusp”) for breaking the DR predictions; however, when approaching a critical line $N_{c,DR}(d)$ (not shown here,
but starting from $N = 18$ for $d = 4 + \epsilon$ and going to $d \approx 5$ when $N = 12$), the critical exponents continuously tend to their DR value. Above the line, the stable (more precisely, once unstable) FP is nonanalytic, but the nonanalyticity is now too weak to break DR. (FPs with a “cusps” can still be found, but they are unstable in several directions and correspond therefore to multicritical points, unreachable from generic initial conditions.)

Finally, in region II we find two zero-temperature nonanalytic FPs: one is attractive and describes a QLRO phase and the other one is once unstable and governs the transition between the disordered and the QLRO phases. The exponent $\eta$ characterizing the power-law decay of the connected correlation function is shown in Fig. 2 for a range of values of $N$. One can see that for $N$ less than a critical value $N_c = 2.83...$ and for $d < 4$, which corresponds to region II, the two FPs coalesce for a value $d_c(N)$ which then determines the lower critical dimension below which no phase transition is observed. The most striking outcome is that, contrary to what is usually believed, no QLRO, i.e., no Bragg glass phase, exists in $d = 3$ for the RFXYM; indeed, $d_c(N = 2) \approx 3.8$. Our study of the RG flow for $d = 3$ and $N = 2$ shows signatures of a “ghost” nontrivial FP, presumably lying on the imaginary axis not too far from the physical plane of coupling constants (which could explain the behavior found in Monte Carlo simulation [10], see Fig. 3, but the flow goes at large distance to the trivial disordered FP. We shall come back to this point later. Note that an estimate of the uncertainty of the present nonperturbative but approximate RG treatment is given by considering the point in Fig. 1a for $d = 2$: in this (probably most unfavorable) case, the theory predicts $N_{lc} \approx 1.15$ instead of the exact result $N_{lc} = 1$. As shown for the pure $O(N)$ model [12], this could be improved by solving the full first order of the derivative expansion.

Another intriguing question raised by this similarity is the possible equivalence between the $(N_c = 2.83..., d = 4)$ point of the RFO(N)M and the $(N_c = 2, d = 2)$ point of the pure model (compare Figs 1a and b). Our nonperturbative FRG solution of the former shows a behavior very reminiscent of the KT transition, but the approximation used is unable to rigorously locate a line of FPs as would occur for a KT transition. As studied in detail for the pure $d = 2$ XY model [13, 15], a line of “almost” FPs is found: this is enough to observe a behavior almost indistinguishable from the KT scenario, but not satisfying to prove the existence of a KT transition. To resolve this matter, we have considered the perturbative FRG analysis around $d = 4$. The 1-loop $\beta$-function (or rather functional) for the renormalized 2-replica potential (to which our nonperturbative treatment exactly reduces near $d = 4$) can be shown to vanish identically in $d = 4$ when $N = N_c = 2.83...$ for arbitrary values of the disorder strength, much like it does in the pure $d = 2$ XY model for arbitrary values of the temperature; but what about the 2-loop $\beta$-function? To answer this question we have calculated (independently from the authors of the recent ref. [18]) the $\beta$-functions of the RFO(N)M around $d = 4$ to 2 loops, starting with the nonlinear-sigma model associated with Eq. 15. The calculation follows the treatment developed for the pure model [16], with the bare action expressed in terms of renormalized nomenclatural RG equations [16] and from known exact results, is shown in Fig. 1h. Regions I, II and II are the exact counterparts of those in Fig. 1h. Indeed, and although never mentioned, QLRO exists in the pure model even aside from the special KT transition for $N = 2$ and $d = 2$. However, it occurs in an unphysical region, $1 < N < 2$, $1 < d < 2$! It remains that the topology of the RFO(N)M phase diagram is very similar to that of the pure model. Again, there is no exact DR-like property, the lower critical dimension $d_c(N)$ being shifted by an $N$-dependent value between the two models.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig2.png}
\caption{Prediction for the $d$-dependence of the anomalous dimension $\eta$ for different values of $N$ (by steps of 0.2). Below $N_c = 2.83...$ and for $d < 4$ two nontrivial FPs are found and coalesce for $d = d_c(N)$ (dots).}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig3.png}
\caption{Inverse correlation length vs disorder strength for the 3-d RFXYM. The crossover around $\Delta = 5$ is the remain of the FPs found above $d_c$. $\Delta$ parametrizes the initial conditions of the RG flow and $1/\xi$ is estimated as in Ref. 14.}
\end{figure}
dimensionless quantities,

\[ S(\{ \pi_a \}) = -\frac{k^d}{2Z_T r^{1/2}} \sum_{x,a,b=1} \left[ R(Z_a \pi_a \cdot \pi_b + \sigma_a \sigma_b) + \delta[R] \right] \]

\[ + \frac{k^{d-2}}{2Z_T} \int_x \sum_{a=1}^n \left[ \frac{Z}{\sqrt{\gamma}} (\partial \sigma_a)^2 + (\partial \sigma_a)^2 - 2h \frac{Z}{\sqrt{\gamma}} \sigma_a \right] \]

where the \((N - 1)\)-component fields \( \pi_a (\pi_a = \frac{Z}{\sqrt{\gamma}} \Pi_a) \) characterize the fluctuations in the directions orthogonal to the external field \( h (h = \frac{Z}{\sqrt{\gamma}} Z_T^{-1/2} H) \) while \( \sigma_a = \sqrt{1 - \frac{Z}{\sqrt{\gamma}} \Pi_a^2} \) is \( k^{d-2} Z_T^{-1} T \) and \( R(z) \) is the renormalized 2-replica potential that we seek to relate to the bare term \( R_0(z) = \Delta z \) via the counterterms \( \delta[R] \), such that \( R_0 = k^{d-d}(R + \delta[R]) \). We then carry out the perturbation expansion around the gaussian theory characterized by the propagator \( G_{ab}^{(0)}(q) = t Z_T k^{2d}(Z_q^2 + Z_T^{-1} Z_T^{-1} H) \) where \( Q_L \) is given in refs \([12, 21]\) and the 2-loop contribution reads

\[ 2 \beta_2[r] = y(yr'' - 3yr'' - r' - 1) - (N - 2) \left[ yr'' - 3yr'' + yr' - 1 \right] \]

\[ + a^2 \left[ (N + 2) yr'' - (3N - 2) yr' + 8K(N - 2)r \right], \]

where \( y = 1 - z^2 \) and \( a = \lim_{(z-1)^{-1}(1 - z^2)^{1/2}} \pi(z) \). One can check that with the change of variable \( z = \cos \phi \) and with \( K = 2a_\gamma \), the above equation is identical to the result recently obtained by Le Doussal and Wiese \([18]\).

As expected, DR is recovered by setting the “anomalous” term \( a \) to zero. The coefficient \( K \) in the above expression is left so far as an unknown: it requires additional calculations (in progress) in order to fully determine the 2-loop \( \beta \)-functions. However, its value is not needed to answer the above question, i.e., to know if for \( N_c = 2.83 \ldots \) and \( d = 4 \) the \( \beta \)-function for \( R(z) \) vanishes identically to two loops for arbitrary disorder strength \( R'(1) \). This latter property is only true if both \( \beta_1[r] \) and \( \beta_2[r] \) vanish identically for the same function \( r(z) \). One can check that this is impossible, irrespective of the value of \( K \). We can thus conclude that the special point \( (N_c = 2.83\ldots,d = 4) \) does not correspond to a KT transition. Actually, as also obtained in our nonperturbative approach (see Fig[1]), the only FP found for the \((N_c, d = 4) \) model is at zero renormalized disorder strength, which allows a perturbative analysis of the two nonanalytic fixed points for \( d < 4 \) and \( N < N_c \) \([18]\).

One further consequence of the absence of a KT transition is that, contrary to what occurs in the pure model near \((N = 2, d = 2) \), the line of lower critical dimension approaches the \((N_c, d = 4) \) point with an infinite slope (compare Figs[11] and [b]). This reinforces our finding that \( d = 3 \) is safely below the lower critical dimension of the RFXYM. Whether or not this conclusion implies that no Bragg glass phase should exist in disordered high-\( T_c \) superconductors is however an open question. The mapping from these latter to the RFXYM is only valid in the low-\( T \) phase where an essentially elastic description can be used. It may well be that the mechanism by which QLRO is destroyed in the RFXYM, a mechanism that does not explicitly invokes the presence of dislocations, is specific to that model.

\[ \]


