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LETTER TO THE EDITOR

Quadratic spatial anisotropy model with singularities: comparison between exact solution and renormalisation group for crossover dependence

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Abstract. We consider a quadratic spatially anisotropic model, with a $(d-d_\perp)$-dimensional defect hyperplane embedded in d space, and provide the exact solution for the defect hyperplane correlation length and susceptibility for arbitrary $\varepsilon_\perp = 2 - d_\perp > 0$. A perturbation expansion in the defect surface interaction displays *singularities* in $\varepsilon_\perp$ and therefore requires renormalisation. Renormalisation group descriptions of the crossover dependence of scaling amplitudes on the strength of the surface interaction are compared with the exact analytic solution. The close analytic similarities between the renormalisation group crossovers for scaling amplitudes in the spatially anisotropic model and in ϕ^4 field theory suggests that our results for the former have a bearing on the expected faithfulness of the latter.

Renormalisation group calculations, especially those using ε expansion techniques (Wilson and Kogut 1974, Fisher 1974, Amit 1978), have been extremely useful in studying critical phenomena. Models have been considered using ϕ^4 (or $\phi^4 - \phi^6$) field theories, possibly including quadratic anisotropy either in spin or in position space (Aharony 1976, Domany *et al* 1977, Blankschtein and Aharony 1983, Goldschmidt 1983, Nemirovsky and Freed 1985a, b). Some studies focus primarily on the computation of critical exponents where there are independent numerical checks on the accuracy of renormalisation group results, but there is also a growing interest in the description of the crossover between various fixed points limits. Complicated crossover analyses are generally limited to the use of low-order perturbation expansions and serious questions, therefore, exist concerning the accuracy of these low-order calculations.

Exact solutions to non-trivial models are immensely useful in providing guidance through the quagmire of approximation techniques and here we provide one such example in which approximate low-order renormalisation group descriptions can be compared with *exact closed form* solutions. More explicitly, we consider the n-component ϕ^2 model with a spatial anisotropy involving the interaction with a penetrable $(d-d_\perp)$-dimensional planar hypersurface. The free energy functional for this model is

$$F = \int \mathrm{d}^d r[\tfrac{1}{2}(\nabla\phi)^2 + \tfrac{1}{2}t\phi^2 + \tfrac{1}{2}c_0\,\delta^{d_\perp}(\boldsymbol{r}_\perp)\phi^2] \qquad (1)$$

where $\phi = (\phi_1, \phi_2, \ldots, \phi_n)$ is the order parameter, $t \propto T - T_0$ is the bare reduced temperature with T_0 the mean-field bulk critical temperature, c_0 is the bare surface interaction parameter, $\boldsymbol{r}_\perp$ is the component of position vectors orthogonal to the surface and $\delta^{d_\perp}$ is a $d_\perp$-dimensional delta function. Equation (1), when supplemented by a

$g_0\phi^4$ contribution and restricted to $d_\perp = 1$, has been extensively used to analyse critical phenomena near an interacting impenetrable surface (Bray and Moore 1977, Diehl and Dietrich 1980, 1981a, b, Nemirovsky and Freed 1985a, b, Binder 1983). The restriction to $d_\perp = 1$ enables the full propagator, evaluated from F of (1), to be computed exactly. The additional ϕ^4 contribution may then be treated perturbatively with ε expansion methods and with the zeroth-order propagator including a full crossover dependence on c_0 (Goldschmidt 1983, Nemirovsky and Freed 1985a).

Here we follow the suggestion of Kosmas (1985) within a ($n \to 0$ limit) polymer context of analytically continuing to continuous $d_\perp$. The perturbation expansion of (1) in powers of c_0 is shown here to contain singularities as a function of $d_\perp$ and therefore to require renormalisation. Burkhardt and Eisenriegler (1981) have previously noted that the model (1) has the critical dimensionality $d_\perp = 2$ at which it acquires logarithmic corrections. This is equivalent to our observation that the model also leads to *singularities* in $\varepsilon_\perp = 2 - d_\perp$ and therefore requires renormalisation. Following Burkhardt and Eisenriegler (1981), we consider this class of models to be non-trivial because they contain singularities and they provide exactly solvable examples against which to test renormalisation group concepts. For instance, a ϕ^2 theory with mass anisotropy (a simple homogeneous quadratic anisotropy) does not produce singularities and is, therefore, a trivial model in this sense.

The perturbation expansion of (1) in powers of c_0 is shown here to be amenable to exact closed form solution by standard renormalisation group $\varepsilon_\perp$ expansions, thereby providing a non-trivial zeroth-order model for critical phenomena at defect planes, etc. The study of this model is also a useful prelude to investigating the double crossover produced by the more general case with both $g_0\phi^4$ and $c_0\delta(\boldsymbol{r}_\perp)\phi^2$. This crossover has been considered for $d_\perp = 1$ (Goldschmidt 1983, Nemirovsky and Freed 1985a) and is rather complicated. There are, however, indications that the approach of taking $d_\perp$ to be a continuous variable leads to essential simplifications.

The renormalisation group crossover in c, the renormalised counterpart of c_0, is compared with the exact solution to assess the faithfulness of the renormalisation group description of crossovers. We are aware of no other such comparisons using non-trivial models. Furthermore, as explained below, although the ε expansions for ϕ^4 field theories are more complicated than the $\varepsilon_\perp$ ones for the anisotropic ϕ^2 models, the formal structure of the renormalisation constants and the Gell-Mann-Low functions is identical to order ε and $\varepsilon_\perp$, respectively (only numerical coefficients differ). Since the solution to the renormalisation group equation depends on the Gell-Mann-Low functions, the structure of this solution is very similar for the anisotropic ϕ^2 models and ϕ^4 models and the resultant crossover behaviours for scaling amplitudes have many features in common.

The two-point correlation function $G(\boldsymbol{k})$ in momentum space is obtained from (1) using standard techniques (Amit 1978) as

$$G(\boldsymbol{k}) = G^0(\boldsymbol{k})\left[1 - c_0(2\pi)^{-d_\perp}\int \frac{\mathrm{d}^{d_\perp}q_\perp}{\boldsymbol{q}_\perp^2 + \boldsymbol{k}_\parallel^2 + t} + \left(c_0(2\pi)^{-d_\perp}\int \frac{\mathrm{d}^{d_\perp}q_\perp}{\boldsymbol{q}_\perp^2 + \boldsymbol{k}_\parallel^2 + t}\right)^2 \cdots\right] \qquad (2)$$

where the free propagator is $G^0(\boldsymbol{k}) = (\boldsymbol{k}^2 + t)^{-1}$, $\boldsymbol{k}_\perp$ is the Fourier conjugate of $\boldsymbol{r}_\perp$ and $\boldsymbol{k} = (\boldsymbol{k}_\perp, \boldsymbol{k}_\parallel)$. The geometric series in (2) is readily summed to give the *exact* closed form solution

$$G(\boldsymbol{k}) = G^0(\boldsymbol{k})\left(1 + c_0(2\pi)^{-d_\perp}\int \frac{\mathrm{d}^{d_\perp}k_\perp}{\boldsymbol{k}_\perp^2 + \boldsymbol{k}_\parallel^2 + t}\right)^{-1}. \qquad (3)$$

The integral in (3) is evaluated using dimensional regularisation:

$$\int \frac{d^{d_\perp}k_\perp}{k_\perp^2+k_\parallel^2+t}=\frac{\pi^{d_\perp/2}\Gamma(\varepsilon_\perp/2)}{(k_\parallel^2+t)^{\varepsilon_\perp/2}} \tag{4}$$

where $\Gamma(\varepsilon_\perp/2)$ is the gamma function. Equation (1) yields a two-point function in (3) which is independent of the number of spin components.

The exact in-plane correlation length $\xi_\parallel^2$ in the hypersurface of dimensions $d_\parallel(=d-d_\perp)$ and the susceptibility χ are evaluated, respectively, as

$$\begin{aligned}\xi_\parallel^2&=-\left(\frac{\partial G(k)}{\partial k_\parallel^2}(G(k))^{-1}\right)_{k=0}\\&=t^{-1}\{1-c_0(4\pi)^{-d_\perp/2}\Gamma(1+\varepsilon_\perp/2)t^{-\varepsilon_\perp/2}[1+c_0(4\pi)^{-d_\perp/2}\Gamma(\varepsilon_\perp/2)t^{-\varepsilon_\perp/2}]^{-1}\}\end{aligned} \tag{5}$$

$$\chi=G(k=0)=t^{-1}[1+c_0(4\pi)^{-d_2/2}\Gamma(\varepsilon_\perp/2)t^{-\varepsilon_\perp/2}]^{-1}. \tag{6}$$

In addition to the divergence of (5) and (6) at the bulk critical point $t=0$, when c_0 is negative, both $\xi_\parallel^2$ and χ diverge at the shifted critical temperature T_c

$$T_c=T_0+[|c_0|(4\pi)^{-d_\perp/2}\Gamma(\varepsilon_\perp/2)]^{2/\varepsilon_\perp} \tag{7}$$

corresponding to the surface transition.

Now we consider the renormalisation group calculation of the crossover dependence of $\xi_\parallel^2$ on the renormalised counterpart of c_0. For this purpose it is convenient to rewrite (5) as the bare perturbative expansion

$$\begin{aligned}\xi_\parallel^2=t^{-1}\{1&-\Gamma(1+\varepsilon_\perp/2)u_s^0(\kappa/t)^{\varepsilon_\perp/2}[1-u_s^0(\kappa/t)^{\varepsilon_\perp/2}\Gamma(\varepsilon_\perp/2)\\&+(u_s^0)^2(\kappa/t)^{\varepsilon_\perp}\Gamma^2(\varepsilon_\perp/2)-\ldots]\}\end{aligned} \tag{8}$$

where the dimensionless surface interaction (anisotropy) parameter is defined as

$$u_s^0=c_0\kappa^{-\varepsilon_\perp/2}(4\pi)^{-d_\perp/2} \tag{8a}$$

with κ having the dimensions of t. Equation (8) clearly displays poles in $\varepsilon_\perp$ as arising from the factors of $\Gamma(\varepsilon_\perp/2)$ when $\varepsilon_\perp\to 0$. Furthermore, the truncated form of (8) is undefined for $c_0\to\infty$. Hence, equation (8) requires standard renormalisation.

As usual, define the renormalisation constant Z by

$$u_s^0=Z_s u_s. \tag{9}$$

Simple algebra shows that

$$Z_s=1+2u_s/\varepsilon_\perp+(2u_s/\varepsilon_\perp)^2+\ldots \tag{10}$$

eliminates poles through second order in $\varepsilon_\perp$ in (8). Then (8) through second order in the renormalised u_s is found as

$$\xi_\parallel^2=t^{-1}\{1-u_s(\kappa/t)^{\varepsilon_\perp/2-u_s}[1+(\varepsilon_\perp/2-u_s)\psi(1)]+O(u_s^3)\} \tag{11}$$

where the gamma functions in (8) have been $\varepsilon_\perp$ expanded and $\psi(1)$ is the psi function. We now provide a brief account of the renormalisation group crossover analysis in terms of a crossover scaling field $\xi_s=\xi_s(u_s)$. Then the approximate renormalisation group result is compared with the exact solution by introducing the relation between the original model parameter c_0 and the renormalised one.

Consider a generic property $F = F(c_0, t)$ (e.g. $\xi_{\parallel}^2$) which is written using (8*a*) and (9) in terms of u_s^0 or u_s, respectively, as

$$F = t^p F_1[(\kappa/t)^{\varepsilon_\perp/2} u_s^0] \tag{12a}$$

$$F = t^p F_1[(\kappa/t)^{\varepsilon_\perp/2} Z_s u_s] \tag{12b}$$

where F is assumed to scale naively with temperature as t^p and F_1 is a dimensionless function. The renormalisation group equation for F_1 is

$$\left(\kappa\frac{\partial}{\partial\kappa} + \beta(u_s)\frac{\partial}{\partial u_s}\right) F_1[(\kappa/t)^{\varepsilon_\perp/2} Z_s u_s] = 0 \tag{13}$$

where the Gell-Mann–Low function $\beta(u_s)$ is defined as usual by

$$\beta(u_s) = \kappa(\partial u_s/\partial\kappa). \tag{14}$$

The perturbative expansion for $\beta(u_s)$ is obtained from (9) and (10) as

$$\beta(u_s) = -(\varepsilon_\perp/2)u_s(1 - 2u_s/\varepsilon_\perp) + O(u_s^4) \tag{15a}$$

giving the non-trivial fixed point to second order in $\varepsilon_\perp$ of

$$u_s^* = \varepsilon_\perp/2 + O(\varepsilon_\perp^3). \tag{15b}$$

The exact closed form results (5), (8*a*) and (9) also enable the proof by direct substitution that $Z_s = (1 - 2u_s/\varepsilon_\perp)^{-1}$ in (13) provides the solution to all orders and, consequently, that $u_s^* = \varepsilon_\perp/2$ is exact through infinite orders in $\varepsilon_\perp$.

The general solution to the renormalisation group equation (13) is

$$F = t^p F_1\left[(\kappa/t)^{\varepsilon_\perp/2} u_s^* \exp\left(-\int_{u_s^*/2}^{u_s} \frac{(\varepsilon_\perp/2)\,\mathrm{d}x}{\beta(x)}\right)\right] \tag{16a}$$

where the lower integration limit has been chosen for convenience. It is natural to introduce the dimensionless crossover scaling field ζ_s such that

$$F = t^p F_2(\zeta_s) \tag{16b}$$

$$\zeta_s = (\kappa/t)^{\varepsilon_\perp/2} \exp\left(-(\varepsilon_\perp/2)\int_{u_s^*/2}^{u_s} [\mathrm{d}x/\beta(x)]\right) \tag{16c}$$

which, upon use of (15*a*), is found to second order as

$$\zeta_s = (\kappa/t)^{\varepsilon_\perp/2}[u_s(1 - u_s/u_s^*)^{-1}](u_s^*)^{-1}. \tag{16d}$$

A comparison of (12*a*) and (12*b*) shows that the scaling field ζ_s of (16) can also be expressed in terms of u_s^0 by

$$\begin{aligned}\zeta_s &= (\kappa/t)^{\varepsilon_\perp/2}(u_s^0/u_s^*)\\ &= [c_0(4\pi)^{-d_\perp/2} t^{-\varepsilon_\perp/2}]/u_s^*\end{aligned} \tag{17}$$

where the second equality uses (8*a*). This relation between ζ_s and c_0 facilitates the comparison between the approximate renormalisation group and the exact solutions.

Conversion of the renormalised expansion (11) into a form consistent with (16*b*) requires the inversion of $\zeta_s[(u_s/u_s^*), \kappa/t]$ from (16*d*) which is

$$u_s/u_s^* = \zeta_s(\kappa/t)^{-\varepsilon_\perp/2)}[1 + \zeta_s(\kappa/t)^{-\varepsilon_\perp/2}]^{-1}. \tag{18}$$

It may be shown that (18) is exact to infinite order in $\varepsilon_\perp$. Substituting (18) for u_s into (11) for $\xi_\parallel^2$ and retaining terms to order $\varepsilon_\perp^2$ provides the renormalisation group crossover scaling form

$$\xi_\parallel^2 = t^{-1}\{-u_s^*[\zeta_s/(1+\zeta_s)][1+u_s^*(1+\zeta_s)^{-1}\psi(1)]+O(\varepsilon_\perp^3)\}. \quad (19)$$

On the other hand, using the relation (17) between c_0 and the scaling field ζ_s the exact solution is represented in the more compact form in terms of ζ_s,

$$\xi_\parallel^2 = t^{-1}\{1-u_s^*\zeta_s\Gamma(1+\varepsilon_\perp/2)[1+\zeta_s\Gamma(1+\varepsilon_\perp/2)]^{-1}\}. \quad (20)$$

It is interesting to note that the second-order renormalisation group crossover expansion (19) is just the $\varepsilon_\perp$ expansion of the full solution (20) truncated to order $\varepsilon_\perp^2$, where the $\varepsilon_\perp$ dependence of the scaling field ζ_s *is not* $\varepsilon_\perp$ *expanded.* The approximate (19) and the exact (20) becomes identical in the limit of $\zeta_s \to \infty$. This situation is very fortunate and it leads to the expectation in general cases that if the exact crossover solution in terms of the scaling field is only a finite power series in ε, the renormalisation group crossover is also only a finite ε expansion. In fact, we show below that this observation is correct within the context of the present model.

Standard ϕ^4 field theory leads to rather complicated ε expansions. However, the crossover dependence in ϕ^4 field theories is based on the solution of the appropriate renormalisation group equation which is found to bear a strong correspondence to that emerging from the model (1). For instance, n-component ϕ^4 field theory yields Z_u and $\beta(u)$ as (Amit 1978) the expansions in dimensionless renormalised u

$$Z_u = 1+(n+8)u/6\varepsilon+[(n+8)^2/36\varepsilon^2-(2n+14)/24\varepsilon]u^2+O(u^3) \quad (21)$$

$$\beta(u) = -\varepsilon u[1-(n+8)u/6\varepsilon+(3n+14)u^2/12\varepsilon]+O(u^4) \quad (22)$$

which are similar in form to (10) and (15a), respectively. The main difference is that (15a) actually terminates at order u_s^2, while (22) does not. The crossover in u involves the scaling field ζ which is of the form of (16c) with an extra factor on the right-hand side arising from 'mass' renormalisation. The net result is that the analogue of (16d) develops ε-dependent corrections to the exponent of $(1-u/u^*)$ and that the natural scaling field η becomes a function of ζ. Thus, a bare scaling amplitude of the form $(1+au_0+bu_0^2)$ has a crossover structure like

$$1+au^*\eta(1+\eta)^{-1}+b(u^*)^2\eta^2(1+\eta)^{-2} \quad (23)$$

where $\eta=\zeta(1+\zeta)^{\varepsilon/(n+8)}+O(\varepsilon^2)$. Equation (23) then has a similar form to (19) for the spatially anisotropic model. While the model (1) is considerably simpler than the general ϕ^4 theory, the close similarities between the analytical structure of the renormalisation group crossovers for scaling amplitudes suggests that our exact results for (1) have a counterpart in ϕ^4 field theory.

As mentioned in the introductory paragraphs, exact solutions are extremely useful in providing guidance in renormalisation group manipulations. For example, our exact solution (20) suggests an alternative (actually an optimal) choice for the definition of the interaction parameter to improve convergence of the $\varepsilon_\perp$ expansion. Modifying the definition (8a) for u_s^0 to

$$\hat{u}_s^0 = u_s^0\Gamma(1+\varepsilon_\perp/2) \quad (24)$$

converts (20) into the exact crossover form

$$\xi_\parallel^2 = t^{-1}[1-\hat{u}_s^*\hat{\zeta}_s/(1+\hat{\zeta}_s)]. \quad (25a)$$

Applying the renormalisation group analysis to the perturbation expansion (8) as written in terms of the variable $\hat{u}_s^0$ yields

$$\xi_\parallel^2 = t^{-1}[1 - \hat{u}_s^* \hat{\zeta}_s/(1+\hat{\zeta}_s) + O(\varepsilon_\perp^3)] \tag{25a}$$

where $\hat{u}_s^* = u_s^* = \varepsilon_\perp/2$. The second-order term is absent in (25*b*); hence, the order ε^2 renormalisation group result (25*b*) *is identical to the exact solution* (25*a*), *entirely because of the modified definition of the interaction parameter* $\hat{u}_s^0$, a definition which is sensible because factors of $\Gamma(1+\varepsilon_\perp/2)$ appear as the coefficient of $c_0 t^{-\varepsilon_\perp/2}$ in each order of perturbation theory (cf (8) with $\Gamma(\varepsilon_\perp/2) = (2/\varepsilon_\perp)\Gamma(1+\varepsilon_\perp/2)$). (The same considerations apply to χ.) This example provides the important lesson that ε-expansion errors may be minimised by absorbing ε-dependent factors into the definition of the coupling constants when some guidance is available from the structure of higher-order contributions or, perhaps, from exact results in special limits.

The above conclusions are not artefacts of dealing with a convergent perturbation expansion as can be seen by considering a hypothetical model which results in the asymptotic Borel summable perturbation expansion for the dimensionless property Q:

$$Q = \left(\sum_{n=0}^{\infty} (-u_0/\varepsilon)^n (\kappa/t)^{\varepsilon n} \Gamma(2+n\varepsilon)\right)\left(\sum_{n=0}^{\infty} (-u_0/\varepsilon)^n (\kappa/t)^{\varepsilon n} \Gamma(1+n\varepsilon)\right)^{-1}. \tag{26}$$

Equation (26), unlike (8), has zero radius of convergence as a power series in u_0 in both numerator and denominator, but the Borel summed form provides a well defined 'exact' result. Application of the renormalisation group method to the perturbation series (26) leads to results very similar to (19). For example, it can be shown that the second-order renormalisation group crossover for Q is again simply the ε expansion to order ε^2 of the exact result (26) where the scaling variable is defined as $\zeta = (\kappa/t)^\varepsilon (u_0/\varepsilon)$ and *is not* ε expanded. This simple mathematical model emphasises our expectation that the ϕ^4 field renormalisation group crossover for scaling amplitudes very likely exhibits a similar approximation to the exact results.

The spatially anisotropic spin model considered here is the Laplace transform of one which is concurrently applied (in the $n \to 0$ limit) to polymers (Douglas *et al* 1986). The polymer case is mathematically more complicated as the exact solutions are represented in terms of Mittag-Lefler functions (Hardy 1949). On the other hand, the polymer model is richer in the sense of also being well defined for any $c_0 < 0$ and the exact solution exhibits a meaningful crossover between $c_0 \to -\infty$ and $c_0 \to +\infty$. The renormalisation group crossover provides a good approximation for $\zeta_s > -1$, but it breaks down for $\zeta_s \leq -1$, providing a useful model for studying potential generalisations of renormalisation group methods to treat problems with negative coupling constants (attractive interactions) where many interesting phenomena occur.

The truism that exact solutions to simple but non-trivial physical models provide important insights into approximate solutions of similar but more complicated problems is exemplified by our comparison of renormalisation group descriptions of the crossover behaviour of a spatially anisotropic ϕ^2 model with that of the exact solution. The crossover for scaling amplitudes is represented in the scaling field ζ_s in the form of $\zeta_s(1+\zeta_s)^{-1}$, whereas crossovers for scaling amplitudes in ϕ^4 field theory also involve expansions in powers of $\varepsilon\eta(1+\eta)^{-1}$ with η the appropriate scaling field. On the basis of the ϕ^2 model we anticipate that the ϕ^4 renormalisation group crossover for scaling amplitudes is likely to be just the ε expansion of the exact solution where the scaling field η is *not expanded in* ε. Monte Carlo simulations of crossovers can, perhaps, test this matter further and thereby lend further support to renormalisation group treatments of crossover behaviour.

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