Critical fluctuations in superconductors

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The field-theoretical gauge model for a superconductor, generalized to a $n/2$-component complex order parameter, is renormalized in two-loop order. The question of a fluctuation-induced first-order transition is discussed, and the crossover functions for the effective exponents and the amplitude ratio of the specific heats above and below T_c are given.

INTRODUCTION

In conventional superconductors the Ginzburg criterion shows that critical fluctuations will be unobservable in the experimentally accessible temperature region near T_c. Especially the nature of the phase transition, whether it is a slightly first-order or a second-order one, cannot be decided. However, the smallness of the background correlation length in the new high-T_c superconductors leads to a much wider critical region, and critical effects may become observable.[1,2] We therefore reconsider fluctuation effects in superconductors in two-loop order, described by an Abelian Higgs model, within a nonasymptotic renormalization-group theory.

Several aspects have already been discussed in the literature. In one-loop order, runaway solutions for the renormalized couplings in an expansion in $\epsilon=4-d$ were found for the number of order-parameter components $n<n_c=365.9$.[3] These have been interpreted as the appearance of a first-order phase transition and were confirmed for $8.94<n<365.9$ by considering the equation of state.[4] On the other hand, results obtained in a lattice model[5] and in an expansion of the corresponding nonlinear σ model[6] in $2+\epsilon$ did not show a first-order phase transition, and arguments were given by Lawrie[7] that low-order perturbation theory fails for $n<8.94$.

Crossover functions, calculated by the matching technique of Nelson and Rudnick[8,9] to first order in ϵ, for the order-parameter susceptibility and the effective exponent γ_{eff} near the fluctuation-induced first-order transition, were discussed by Chen *et al.*,[10] and an expression for the crossover function of the specific heat was given.

Here we perform a renormalization of the theory in *two-loop order,* and calculate the β functions for the flow of the renormalized coupling parameter (fourth-order coupling of the order parameter and coupling to the gauge field) and the ζ functions appearing in the vertex functions and determining the critical exponents. For the special case $n=2$ the β functions were already calculated by Vladimirov and Shirkov;[11] see also van Damme.[12] The strict ϵ expansion is discussed, and a new lower value of n_c appears below which the fixed point value of the renormalized charge squared (e^2) is negative. In addition, the fixed point value of the fourth-order coupling u^* remains complex below $n_c=365.9$. On the other hand, a large reduction of the borderline value to $n_c=2.4$ is found from the two-loop-flow β functions without expansion of the solution in ϵ. We further discuss the nonuniversal crossover functions for the effective exponent of the specific heat and the amplitude ratio of the specific heat above and below T_c.

FLOW EQUATIONS

The system is described by the gauge-invariant Hamiltonian

$$\mathcal{H}=\int d^3x\left[\tfrac{1}{2}t_0|\Psi_0|^2+\tfrac{1}{2}|(\nabla-ie_0\mathbf{A}_0)\Psi_0|^2+\frac{u_0}{4!}\left|\Psi_0\right|^4+\tfrac{1}{2}(\nabla\times\mathbf{A}_0)^2\right], \quad (1)$$

where the complex scalar field Ψ_0 with $n/2$ components is coupled minimally to the vector field $\mathbf{A}_0$. The generalization from the three-dimensional space to d dimensions is as usual. All calculations have been performed in the superconducting (transversal or Landau) gauge. For the coupling constant $e_0=0$ no magnetic fluctuations are induced, and the model reduces to the usual field theory describing a second-order phase transition. The parameter t_0 changes sign at some temperature, and the fourth-order coupling u_0 is to be taken as temperature independent. In order to describe the critical behavior of this model we use the field-theoretical approach, and a gauge-invariant procedure has to be chosen.[13] We adopt the dimensional regularization and the minimal subtraction scheme.[14] For a recent explanation of the method see Ref. 15 and references therein. We introduce renormalized fields and couplings,

$$\Psi_0=Z_\Psi^{1/2}\Psi,\quad \mathbf{A}_0=Z_A^{1/2}\mathbf{A},\quad t_0-t_{0c}=Z_tZ_\Psi^{-1}t,$$
$$e_0^2=Z_e^1Z_A^{-1}Z_\Psi^{-1}e^2\mu^\epsilon S_d^{-1},\quad u_0=Z_uZ_\Psi^{-2}u\mu^\epsilon S_d^{-1},$$

where μ is a reference wave number, t_{0c} a shift, which for the results considered in this paper can be set to zero, and the geometrical factor S_d equals $2^{1-d}\pi^{-d/2}/\Gamma(d/2)$. The Z factors are determined by the condition that all poles at $\epsilon=0$ are removed from the renormalized vertex functions. From a Ward identity we have $Z_\Psi=Z_e$, and the remaining Z factors are found from the vertex functions $\Gamma^{(2,0)}$, $\Gamma^{(0,2)}$, and $\Gamma^{(4,0)}$ (the first index is the number of Ψ fields, the second index the number of A fields).

Since the vertex field is massless, the renormalization has been performed at finite wave vector. The considerable number of two-loop diagrams is reduced by choosing the transverse gauge. A reduction of the computational effort could also be achieved by making use of the formulas given in Refs. 16 and 17. Integrals of pure powers appearing in the calculation are equal to zero in dimensional regularization.[18,19] The result in two-loop order is

$$Z_\Psi = 1 + \frac{1}{\epsilon}\{3e^2 - u^2(n+2)/144 + e^4[(n+18)/4\epsilon - (11n+18)/48]\} , \tag{2}$$

$$Z_A = 1 + \frac{1}{\epsilon}\{-ne^2/6 - ne^4/2\} , \tag{3}$$

$$Z_t = 1 + \frac{1}{\epsilon}\{(n+2)u/6 + u^2[(n+2)(n+5)/36\epsilon - (n+2)/24] + ue^2[-(n+2)(1/2\epsilon - \tfrac{1}{3}) + e^4[(3n+6)/2\epsilon + (5n+1)/4]\} , \tag{4}$$

$$Z_u = 1 + \frac{1}{\epsilon}\{(n+8)u/6 + 18e^4/u + u^2[(n+8)^2/36\epsilon - (5n+22)/36] + ue^2[-(n+8)/2\epsilon + (n+5)/3] + e^4[(3n+24)/\epsilon + (5n+13)/2] + e^6/u[3(n+18)/\epsilon - 7n/2 - 45]\} . \tag{5}$$

The one-loop-order result is in agreement with Refs. 20 and 4. Following the standard procedure we obtain the β functions and the flow equations for the renormalized couplings ($f = e^2$),

$$l\frac{df}{dl} = -\epsilon f + n/6f^2 + nf^3 = \beta_f , \tag{6}$$

$$l\frac{du}{dl} = -\epsilon u + (n+8)/6u^2 - (3n+14)/12u^3 - 6uf + 18f^2 + (2n+10)/3u^2 f + (71n+174)/12uf^2 - (7n+90)f^3 = \beta_u . \tag{7}$$

For $n=2$ the β functions coincide with the results of Refs. 11 and 12. The fixed points of the one-loop part were already discussed extensively in Refs. 3, 10, and 4; we note, however, some new and unusual properties of the second-order ϵ expansion in this model.

In one-loop order, the stable fixed point value for $u^z = u_1\epsilon$ $[f^* = (6/n)\epsilon]$ in a strict ϵ expansion was found to be complex for $n < n_c = 365.9$, namely (Ref. 3),

$$u_1 = \frac{3}{n+8}(1 + 36/n + \sqrt{\sigma}) \tag{8}$$

with

$$\sigma = (1+36/n)^2 - 432(n+8)/n^2$$

being zero at $n = n_c$. The stability exponent λ_u given by $\lambda_u = \partial\beta/\partial u|_{u^*}$ (valid up to two-loop order) turns out to be

$$\lambda_u = -\sqrt{\sigma}\epsilon . \tag{9}$$

This leads to an oscillatory flow in u in one-loop order below n_c, with the solution[10] ($s = |\sigma|^{1/2}$)

$$f(l) = fl^{-\epsilon}/[1 + nf(l^{-\epsilon} - 1)/6\epsilon] , \tag{10}$$

$$u(l) = f(l)\frac{n}{2(n+8)}\left\{ s\tan\left[\frac{s}{2}\ln[f(l)f^{-1}l^{\epsilon}] + \arctan\left[\frac{2(n+8)}{sn}\frac{u}{f} + \frac{n+36}{ns}\right]\right] - \frac{n+36}{n}\right\} \tag{11}$$

with f and u the initial parameters at $l=1$. The separatrix connecting the unstable fixed point ($f^*=0$, $u^* \neq 0$) and ($f^* \neq 0$, $|u^*| = \infty$) is given by

$$u(f) = \frac{n}{2(n+8)}f\left[s\tan\left[\frac{s}{2}\ln\frac{6\epsilon - nf}{6\epsilon} + \frac{\pi}{2}\right] + \frac{n+36}{n}\right] . \tag{12}$$

Only for $f(l=1)$ and $u(l=1)$ below the first branch of $u(f)$ does the flow reach the limit of small f and u values for $l \to \infty$; for all other initial conditions u goes to plus or minus infinity for finite l. Later on we shall only consider flows with the property $f \to 0$ and $u \to 0$ for $l \to \infty$, since otherwise the Gaussian regime is not reached. The fixed

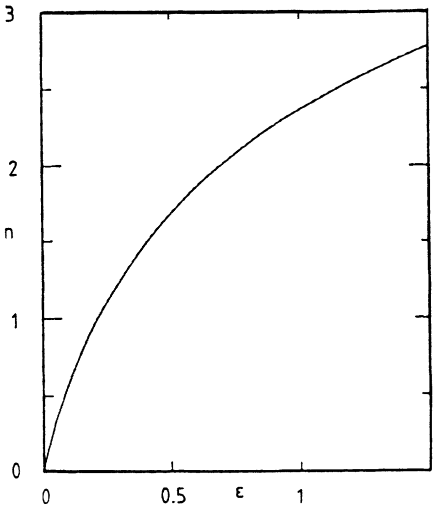

FIG. 1. Borderline in n,ϵ space ($\epsilon=4-d$) between the region where a stable real fixed point exists (above the line) and the region where no stable real fixed point exists (below the line). The dot marks the position of a superconductor.

point in two-loop order is obtained by inserting $u^*=u_1\epsilon+u_2\epsilon^2$ and the two-loop fixed point value for f,

$$f^*=(e^*)^2=(6/n)\epsilon\{1-(36/n)\epsilon\} , \tag{13}$$

into β_u. Then one finds for $n\neq n_c$

$$u_2=A(n)/\sqrt{\sigma}+B(n) , \tag{14}$$

where $A(n)$ and $B(n)$ are real. For $n<n_c$, u^* remains complex; however, a new type of borderline is defined by the condition that the fixed point value, Eq. (13), of $f=e^2$ has to be positive. This is the case at $\epsilon=1$ for $n>36$, whereas the value of f^* without expansion remains positive for all n. One may therefore consider this as an artifact of the expansion procedure.

Instead of a strict ϵ expansion, we now look at the two-loop flow itself given by Eqs. (6) and (7). Thereby we consider the β functions as asymptotic series in the couplings u and f. In fact, one has to use nonperturbative methods like Borel summation[21] for the β functions. This is seen from the usual model with $f\equiv 0$, which has no real stable fixed point for u in two-loop order. However, such a Borel summed form is not available for general n. Therefore, we have to base our discussion on the result of Eqs. (6) and (7). The n-ϵ space is divided into two parts. In one part a real stable fixed point with $u^*>0$ and $f^*>0$ exists; in the other part the stable fixed point is complex. The borderline $n(\epsilon)$ between these two regions is shown in Fig. 1. One observes a drastic reduction of $n(\epsilon)$ at $\epsilon=1$ compared to the one-loop flow and the result of the ϵ expansion. We now turn to the special case $n=2$. We see that in two-loop order [see Fig. 2(a)] the flow is changed considerably from the one-loop-order result.[10,4] It is only in the small region of $f<0.05$ and small u that both flows approximately coincide. This indicates a much worse representation by the low-order result for the β functions than expected from the Ψ^4 theory.

Results for the Borel summed β functions are available for the uncharged model,[22] and we shall use those for $n=2$ in order to give a reliable flow at least for small f values. This correction then leads to the flow shown in Fig. 2(b). As a consequence of this "improvement" a real stable fixed point for $f>0$ and $u>0$ appears even in this case. However, a tiny region exists, for $f<18/(7n+90)$ and very small u, where the flow can escape into the region $f>0$ and $u<0$. New features may arise when the flow equation for u couples to the equation for f, which we expect in three-loop order.

EFFECTIVE EXPONENTS AND AMPLITUDE RATIOS

The critical exponents at a second-order phase transition are determined by the fixed point values of the ζ functions calculated from the renormalization factors Z_i by $\zeta_i=\mu\partial\ln Z_i/\partial\mu$, where the derivative is taken at fixed unrenormalized couplings. Then we get from Eqs. (2)–(4)

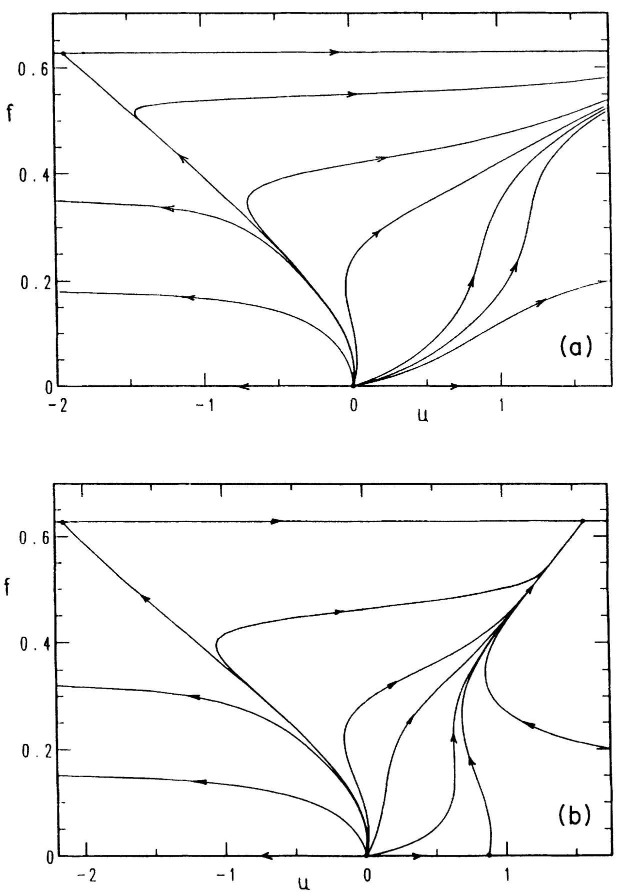

FIG. 2. (a) Flow for the case $n=2$ given by Eqs. (6) and (7) with $\epsilon=1$. No stable real fixed point exists and also the fixed point on the u axis is absent. This should be compared with the one-loop flow given in Ref. 10. (b) Flow for the case $n=2$ given by Eq. (6) and a modified Eq. (7). The f-independent terms in β_u have been replaced by the Borel summed ones given in Ref. 22. This procedure recovers the fixed point on the u axis and also leads to a new stable real fixed point for $f\neq 0$.

$$\zeta_\Psi=-3f+(n+2)/72u^2+(11n+18)/24f^2\ , \quad (15)$$

$$\zeta_t=-(n+2)/6u+(n+2)/12u^2 -2(n+2)/3uf-(5n+1)/2f^2\ , \quad (16)$$

$$\zeta_A=n/6f+nf^2\ . \quad (17)$$

If there exists a stable fixed point, the critical exponent ν of the correlation length and the penetration length (the renormalized Ginzburg parameter reaches a finite fixed point value and has no asymptotic critical temperature dependence), the critical exponent γ of the order parameter susceptibility, and the critical exponent α of the specific heat are given by ($\zeta_\nu=\zeta_\Psi-\zeta_t$),

$$\nu=(2-\zeta_\nu^*)^{-1}\ ,$$
$$\gamma=(2-\zeta_\nu^*)^{-1}(2-\zeta_\Psi^*)\ ,$$
$$\alpha=(\epsilon-2\zeta_\nu^*)(2-\zeta_\nu^*)^{-1}\ .$$

These exponents are different from the values known from the uncharged model, i.e., they are not given by the ^{4}He values as is sometimes stated (see, e.g., Ref. 1). One may consider these asymptotic exponents for the case of the two-loop fixed points; however, because of the large fixed point values and the bad expansion property of the low-order perturbation theory, one exceeds the value of 2 for ζ_ν, which leads to an unphysical negative value of ν at a second-order transition.

Since (i) the asymptotic region cannot be treated theoretically for the reasons already mentioned and (ii) the experimentally accessible regime lies in the precritical region further away from T_c, we discuss several nonasymptotic measurable quantities, such as effective exponents and amplitude ratios.[23] They are calculated from the solution of the renormalization-group equations for the corresponding correlation or vertex function.

One most interesting measurable quantity is the specific heat. The renormalized specific heat $C^\pm$ above (+) and below (−) T_c obeys the following equation:

$$\left[\mu\frac{\partial}{\partial\mu}+\beta_u\frac{\partial}{\partial u}+\beta_f\frac{\partial}{\partial f}+\zeta_\nu\left[2+t\frac{\partial}{\partial t}\right]\right]C^\pm(t,u,f,\mu)=\mu^{-\epsilon}B(u,f)\ , \quad (18)$$

where B comes from the additive renormalization. The formal solution reads

$$C^\pm(t,u,f,\mu)=\mu^{-\epsilon}\exp\left[-\int_1^l[\epsilon-2\zeta_\nu(x)]\frac{dx}{x}\right]\left\{F^\pm(l)-\int_1^l\frac{dy}{y}B(y)\exp\left[-\int_l^y\left[\epsilon-2\zeta_\nu(x)\right]\frac{dx}{x}\right]\right\}\ , \quad (19)$$

where l is an arbitrary flow parameter to be chosen suitably. Usually one chooses $t(l)/l^2=c$, where the constant c may be different above and below T_c. This choice leads to a connection between the temperatures $t^\pm=|T^\pm-T_c|/T_c$ and the flow parameter $l=l(t^\pm)$ via the solution of the flow equation for $t(l)$ with $t(1)\sim t^\pm$,

$$l\frac{dt(l)}{dl}=\zeta_\nu t(l)\ .$$

The scaling functions $F^\pm$ and B have to be calculated by perturbation theory. Since the singular part of the specific heat is defined only up to a constant, we choose this constant such that in the background the specific heat goes to zero. We then define the effective exponent $\alpha_{\text{eff}}^\pm(t_0)$ by

$$C_{0\text{fluct}}^\pm(t)=C_0^\pm(t)-C_0^\pm(t=\infty),\ \alpha_{\text{eff}}^\pm=-d\ln C_{0\text{fluct}}^\pm/d\ln t\ .$$

This exponent is equal to its renormalized counterpart and performing the derivatives one gets $[l=l(t^\pm)]$

$$\alpha_{\text{eff}}^\pm(t^\pm)=\frac{\{F^\pm(l)[\epsilon-2\zeta_\nu(l)]+B(l)-l dF^\pm/dl\}/[2-\zeta_\nu(l)]}{F^\pm(l)+\int_l^\infty\frac{dy}{y}B(y)\exp\left\{\int_l^y[2\zeta_\nu(x)-\epsilon]\frac{dx}{x}\right\}-G(l)} \quad (20)$$

with

$$G(l)=l^\epsilon\exp\left[\int_l^\infty 2\zeta_\nu(x)\frac{dx}{x}\right]\lim_{y\to\infty}[F^\pm(y)y^{-\epsilon}]\ .$$

The scaling function $F^\pm(l)$ depends on l via the couplings $u(l)$ and $f(l)$. Formally the expression is equally valid above and below T_c; the difference is denoted by an index. In the asymptotic limit $l\to 0$, both $\alpha_{\text{eff}}^\pm$ take the asymptotic value α if there is a fixed point; in the background limit $l\to\infty$ we obtain the Gaussian value $\alpha_{\text{eff}}^\pm=\frac{1}{2}$ for $d=3$. The expression for α_{eff} given in paper by Chen *et al.*[10] does not reach this limit because these authors did not subtract the background specific heat. An inspection of the effective exponents with the two-loop flow shows an enhancement for $f\neq 0$ with respect to the case $f=0$. Because of the factor $[2-\zeta_\nu(l)]^{-1}$ the exponents tend to diverge when the ζ_ν function reaches 2. Similar remarks are valid for the appropriately defined γ_{eff}. We wish to note at this point that one has to be careful in the application of asymptotic scaling laws to effective exponents, since these scaling laws are not fulfilled outside the asymptotic region. The nonuniversal deviations may be small, but this has to be verified by experiment.

We may also define an effective amplitude ratio by the ratio of the l-dependent specific heat above and below T_c at the *same* l value.[23] Because of the different dependence of l on the temperature above and below T_c, t^- is then a function of t^+ and the measured specific heats have to be taken at different temperatures above and below T_c:

$$\frac{C^+_{0\,\mathrm{fluct}}(t^+)}{C^-_{0\,\mathrm{fluct}}[t^-(t^+)]} = A[l(t^+)] = \frac{\exp\left[\int_1^l 2\zeta_\nu(l')\frac{dl'}{l'}\right] F^+(l)l^{-\epsilon} - D\lim_{y\to\infty} F^+(y)y^{-\epsilon} - \int_l^\infty dy\, y^{-\epsilon-1}B(y)\exp\left[\int_1^y 2\zeta_\nu(x)\frac{dx}{x}\right]}{\exp\left[\int_1^l 2\zeta_\nu(l')\frac{dl'}{l'}\right] F^-(l)l^{-\epsilon} - D\lim_{y\to\infty} F^-(y)y^{-\epsilon} - \int_l^\infty dy\, y^{-\epsilon-1}B(y)\exp\left[\int_1^y 2\zeta_\nu(x)\frac{dx}{x}\right]} \tag{21}$$

and

$$D = \exp\left[\int_1^\infty 2\zeta_\nu(x)\frac{dx}{x}\right].$$

In the asymptotic region this expression reduces to universal value (for $\epsilon=1$),[23] with appropriate expressions for F^*, B^*, and ζ_ν^*,

$$A^* = \frac{F^{+*} - B^*/(2\zeta_\nu^* - 1)}{F^{-*} - B^*/(2\zeta_\nu^* - 1)}, \tag{22}$$

if there exists a fixed point. The background limit is less easy to obtain. One has to use the one-loop expression for ζ_ν and the one-loop solution of the (renormalization-group) equations [Eqs. (10) and (11)]. Then one has to expand the solution $u^{-1}(l)$

$$u^{-1}(l) = U_1 l^\epsilon + U_2 + O(l^{-\epsilon}). \tag{23}$$

Inserting this into Eq. (21) we find (for $\epsilon=1$)

$$A(\infty) = \frac{F^+(\infty) + B(\infty)}{F_1^-(\infty) + B(\infty) + 2 - n/2 + g(f,u)}. \tag{24}$$

We then have to take the one-loop expressions for the scaling function of the specific heat $F_1^-(l)$, which is defined by $F^-(l) = 3/u(l) + F_1^-(l)$. Contrary to the case where $f=0$ and $A(\infty)=n/4$, one finds for $f\neq 0$ a dependence on the renormalized background values of the couplings both in $F_1^-(\infty)$ and in $g(f,u)$. The contribution $g(f,u)$ comes from the flow as well as from the ζ_ν function, which reads

$$g(f,u) = 3[U_2 - (n+8)/3 + 6U_1 f/(1 - nf/6)]. \tag{25}$$

The one-loop expressions are (i) $B(l)=n/2$, (ii) $F^+(l)=0$, and

$$F_1^-(l) = \left[\frac{6f(l)}{u(l)}\right]^2 \left[\left[\frac{u(l)}{3f(l)}\right]^{1/2} - 1\right],$$

where the scaling functions have been calculated in dimensional regularization in $d=3$ without ϵ expansion, and where the arbitrary parameter l has been chosen by the relations $t(l)/(\mu l)^2 = (\pi/4)^2$ above and $2|t(l)|/(\mu l)^2 = (\pi/4)^2$ below T_c. This leads to the simple connection of the temperature distances in the specific heat $t^- = t^*/2$.[23] Our result for the amplitude ratio is different from the result found in the Gaussian model

$$A(\infty) = n/[4 + (12f_0/u_0)^{3/2}], \tag{26}$$

since in that case the unrenormalized coupling constants are used, whereas we have calculated the amplitude ratio within the renormalized theory. This was necessary because we considered the entire crossover region between the asymptotic critical behavior and the precritical Gaussian region. From Eq. (21) one can also calculate by expansion the deviations from the value given in Eq. (24) in the precritical region. For small values of f_0/u_0 in Eq. (26) or f/u in Eq. (24) (we assume f/u to be small also for large values of the Ginzburg parameter) both corrections to the usual value $n/4$ are very small.

CONCLUSION

We have renormalized the field-theoretic model for a "superconductor" with $n/2$ component complex order parameter. A qualitatively different flow in two-loop order (with respect to one-loop order) for small n is found. No real fixed point exists in two-loop order at $d=3$ and $n=2$; however, for $n>2.36$ a real stable fixed point exists. High-order perturbation theory and Borel summation[15] or other nonperturbative methods are necessary to decide the question of a first-order phase transition and for quantitative calculations of effective exponents or amplitude ratios from the expressions given in our paper.

For an application of this model to real systems further important features should be considered. One may include anisotropies[24,25] or study effects of a dimensional crossover from two to three dimensions. From the experimental side it would be interesting to see if one could find deviations in the critical exponents and amplitude ratios from the Gaussian-model values.

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