

The Effective Potential for Composite Operator in the Scalar Model at Finite Temperature

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We discuss the φ^4 and φ^6 theory defined in a flat D -dimensional space-time. We assume that the system is in equilibrium with a thermal bath at temperature β^{-1} . To obtain non-perturbative result, the $1/N$ expansion is used. The method of the composite operator (CJT) for summing a large set of Feynman graphs, is developed for the finite temperature system. The resummed effective potential and the analysis of the $D = 3$ and $D = 4$ cases are given.

1 Introduction

The conventional perturbation theory in the coupling constant or in \hbar i.e. , the loop expansion can only be used for the study of small quantum corrections to classical results. When discussing quantum mechanical effects to any given order in such an expansion, one is not usually able to justify the neglect of yet higher order. In other words, for theories with a large N dimensional internal symmetry group, there exist another perturbation scheme, the $1/N$ expansion, which circumvents this criticism. Each term in the $1/N$ expansion contains an infinite subset of terms of the loop expansion. The $1/N$ expansion has the nice property that the leading-order quantum corrections are of the same order as the classical quantities. Consequently, the leading order which adequately characterizes the theory in the large N limit preserves much of the nonlinear structure of the full theory. In the next section we derive the effective action to leading order in $1/N$ in D -dimensional space-time and consequently the effective potential. It is known that, in $D > 4$, such theories with φ^4 interaction are in fact free field theory, while in $D < 4$ they have a non-trivial continuum limit as an interacting field theory. For $D = 3$ it has been shown that, in the large N limit, the φ^6 theory has a UV fixed point and therefore must have a second IR fixed point [1]. At least for large N the $(\varphi^6)_{D=3}$ theory is known to be qualitatively different from $(\varphi^4)_{D=4}$ theory.

In a previous work [2] by use of the composite operator formalism and, we re-examine the behavior at finite temperature of the $O(N)$ -symmetric $\lambda\varphi^4$ model in a generic D -dimensional Euclidean space. In the cases

$D = 3$ and $D = 4$, an analysis of the thermal behavior of the renormalized squared mass and coupling constant are done for all temperatures. It results that the thermal renormalized squared mass is positive and increases monotonically with the temperature. It is interesting to stress that the behavior of the thermal coupling constant is quite different in odd or even dimensional space. In $D = 3$, the thermal coupling constant decreases up to a minimum value different from zero and then grows up monotonically as the temperature increases. In the case $D = 4$, it is found that the thermal renormalized coupling constant tends in the high temperature limit to a constant asymptotic value. Also for general D -dimensional Euclidean space, we are able to obtain a formula for the critical temperature of the second order phase transition. This formula agrees with previous known results at $D = 3$ and $D = 4$ [3, 4].

It is well known that the introduction of the φ^6 term generated a rich phase diagram, with the possibility of second order, first order phase transitions or even both transitions occurring simultaneously. This situation defines the tricritical phenomenon. Some systems such antiferromagnets in the presence of a strong external field or the $H\epsilon^3 - H\epsilon^4$ mixture exhibits such behavior. In a previous paper the massive $(\varphi^6)_{D=3}$ model was analyzed at finite temperature at the two-loops approximation. We demonstrate the existence of the tricritical point [5]. A natural extension of this paper was done in ref. [7]. In this paper we proved the existence of the tricritical point using a non-perturbative approach. This was done using the CJT formalism i.e. the composite operator formalism [6].

Here we continue to study the composite operator method. Still studying the φ^6 theory in the large N expansion, the effective potential at finite temperature is calculated. The organization of the paper is the following. In section II we derive the effective potential using the composite operator (CJT) formalism. In section III the thermal effective potential is found for a D -dimensional generic space. Conclusions are given in section IV. In this paper we use $\hbar = c = k_b = 1$.

2 The effective potential (The CJT formalism)

We are interested here in the most general renormalizable scalar field model $\lambda\varphi^4 + \eta\varphi^6$ possessing an internal symmetry $O(N)$, in a generic D -dimensional space-time. Of course, for $D = 4$ this theory is non-renormalizable. In this case to ensure renormalisability we must make $\eta = 0$. Let us define at the beginning the field in a generic D -dimensional space-time. For simplicity we will call this theory a φ^6 model.

Using the method of composite operator developed by Cornwall, Jackiw and Tomboulis [6, 8], Townsend derived the effective potential of φ^6 theory in the $1/N$ expansion for $D = 3$ at zero temperature [9]. This author proved that $1/N$ expansion is consistent for φ^6 to

leading order.

The Lagrangian density of the $O(N)$ symmetry φ^6 theory is :

$$\mathcal{L}(\varphi) = \frac{1}{2}(\partial_\mu\varphi)^2 - \frac{1}{2}m_0^2\varphi^2 - \frac{\lambda_0}{4!N}\varphi^4 - \frac{\eta_0}{6!N^2}\varphi^6, \quad (1)$$

where the quantum field is an N -component vector in the N -dimensional internal symmetry space. For definiteness, we work at zero-temperature; however, the finite temperature generalizations can be easily obtained [10]. We are interested in the effective action $\Gamma(\phi)$ which governs the behavior of the expectation values $\varphi_a(x)$ of the quantum field where ϕ is given by

$$\phi(x) \equiv \frac{\delta W(J)}{\delta J(x)} = \langle 0|\varphi(x)|0 \rangle, \quad (2)$$

where $W(J)$ is the generating functional for connected Green's functions.

$\Gamma(\phi)$ can be shown to be the sum of one-particle irreducible (1PI) Feynman graphs with a factor $\phi_a(x)$ on the external line. We make use of the formalism of composite operator which reduces the problem to summing two particle irreducible (2PI) Feynman graphs by defining a generalized effective action $\Gamma(\phi, G)$ which is a functional not only of $\phi_a(x)$, but also of the expectation values $G_{ab}(x, y)$ of the time ordered product of quantum fields $\langle 0|T(\varphi(x)\varphi(y))|0 \rangle$, i.e.

$$\Gamma(\phi, G) = I(\phi) + \frac{i}{2}TrLnG^{-1} + \frac{i}{2}TrD^{-1}(\phi)G + \Gamma_2(\phi, G) + \dots, \quad (3)$$

where $I(\phi) = \int dx^D \mathcal{L}(\phi)$, G and D are matrices in both the functional and the internal space whose elements are $G_{ab}(x, y)$, $D_{ab}(\phi; x, y)$ respectively and D is defined by

$$iD^{-1} = \frac{\delta^2 I(\phi)}{\delta\phi(x)\delta\phi(y)}. \quad (4)$$

The quantity $\Gamma_2(\phi, G)$ is computed as follows. In the classical action $I(\varphi)$ we have to shift the field φ by

ϕ . The new action $I(\varphi + \phi)$ possesses terms cubic and higher in φ . This defines an interaction part $I_{int}(\varphi, \phi)$ where the vertices depend on ϕ . $\Gamma_2(\phi, G)$ is given by sum of all (2PI) vacuum graphs in a theory with vertices determined by $I_{int}(\varphi, \phi)$ and the propagators set equal to $G(x, y)$. The trace and logarithm in eq.(3) are functional. After these procedures the interaction Lagrangian density becomes

$$\begin{aligned} \mathcal{L}_{int}(\varphi, \phi) = & -\frac{1}{2} \left(\frac{\lambda_0\phi_a}{3N} + \frac{\eta_0\phi^2\phi_a}{30N^2} \right) \varphi_a\varphi^2 - \left(\frac{8\eta_0\phi_a\phi_b\phi_c}{6N^2} \right) \varphi_a\varphi_b\varphi_c - \frac{1}{4!N} \left(\lambda_0 + \frac{\eta_0\phi^2}{10N} \right) \varphi^4 \\ & - \left(\frac{12\eta_0\phi_a\phi_b}{6!N^2} \right) \varphi_a\varphi_b\varphi^2 - \frac{1}{5!} \left(\frac{\eta_0\phi_a}{N^2} \varphi_a\varphi^4 \right) - \frac{\eta_0}{6!N^2} \varphi^6. \end{aligned} \quad (5)$$

The effective action $\Gamma(\phi)$ is found by solving for $G_{ab}(x, y)$ the equation

$$\frac{\delta\Gamma(\phi, G)}{\delta G_{ab}(x, y)} = 0, \quad (6)$$

and substituting the solution in the generalized effective action $\Gamma(\phi, G)$.

The vertices in the above equation contain factors of $1/N$ or $1/N^2$, but a φ loop gives a factor of N provided the $O(N)$ isospin flows around it alone and not into another part of the graph. We usually call such loops bubbles. Then at leading order in $1/N$, the vacuum graphs are bubble trees with two or three bubbles at each vertex. The (2PI) graphs are shown in figure.(1).

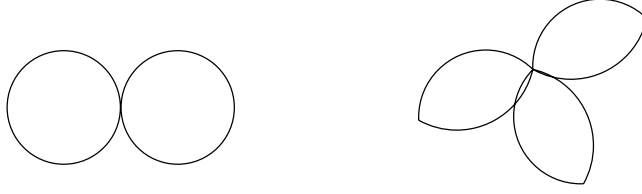


Figure 1.: The 2PI vacuum graphs

It is straightforward to obtain

$$\Gamma_2(\phi, G) = \frac{-1}{4!N} \int d^D x \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) [G_{aa}(x, x)]^2 - \frac{\eta_0}{6!N^2} \int d^D x [G_{aa}(x, x)]^3. \quad (7)$$

Therefore eq.(6) becomes

$$\begin{aligned} \frac{\delta\Gamma(\phi, G)}{\delta G_{ab}(x, y)} &= \frac{1}{2}(G^{-1})_{ab}(x, y) + \frac{i}{2}D^{-1}(\phi) - \frac{1}{12N} \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) [\delta_{ab}G_{cc}(x, x)]\delta^D(x - y) \\ &\quad - \frac{3\eta_0}{6!N} \delta_{ab}[G_{cc}(x, x)]^2\delta^D(x - y) = 0. \end{aligned} \quad (8)$$

Rewriting this equation, we obtain the gap equation

$$\begin{aligned} (G^{-1})_{ab}(x, y) &= D_{ab}^{-1}(\phi; x, y) + \frac{i}{6N} \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) [\delta_{ab}G_{cc}(x, x)]\delta^D(x - y) + \\ &\quad \frac{i\eta_0}{5!N^2} \delta_{ab}[G_{cc}(x, x)]^2\delta^D(x - y). \end{aligned} \quad (9)$$

Hence

$$\frac{i}{2}Tr D^{-1}G = \frac{1}{12N} \int d^D x \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) [G_{aa}(x, x)]^2 + \frac{3\eta_0}{6!N^2} \int d^D x [G_{aa}(x, x)]^3 + cte. \quad (10)$$

Using eqs.(9) and eq.(10) in eq.(7) we find the effective

action

$$\Gamma(\phi) = I(\phi) + \frac{i}{2} \text{Tr}[LnG^{-1}] + \frac{1}{4!N} \int d^D x \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) [G_{aa}(x, x)]^2 + \frac{2\eta_0}{6!N^2} \int d^D x [G_{aa}(x, x)]^3, \quad (11)$$

where G_{ab} is given implicitly by eq.(9). The trace in (11) are both the functional and the internal space. The last two terms on the r.h.s of eq.(11) are the leading contribution to the effective action in the $1/N$ expansion. As usual we may simplify the situation by separating G_{ab} into transverse and longitudinal components, so

$$G_{ab} = \left(\delta_{ab} - \frac{\phi_a \phi_b}{\phi^2} \right) g + \frac{\phi_a \phi_b}{\phi^2} \tilde{g}, \quad (12)$$

in this form we can invert G_{ab} ,

$$(G)_{ab}^{-1} = \left(\delta_{ab} - \frac{\phi_a \phi_b}{\phi^2} \right) g^{-1} + \frac{\phi_a \phi_b}{\phi^2} \tilde{g}^{-1}. \quad (13)$$

Now we can take the trace with respect to the indices of the internal space,

$$G_{aa} = Ng + O(1), \quad (G)_{aa}^{-1} = Ng^{-1} + O(1). \quad (14)$$

From this equation at leading order in $1/N$, G_{ab} is diagonal in a, b . Substituting eq.(14) into eq.(11) and eq.(9) and keeping only the leading order one finds that the daisy and superdaisy resummed effective potential for the φ^6 theory is given by:

$$\Gamma(\phi) = I(\phi) + \frac{iN}{2} \text{tr}(\ln g^{-1}) + \frac{N}{4!} \int d^D x \left(\lambda_0 + \frac{\eta_0 \phi^2}{10N} \right) g^2(x, x) + \frac{2N\eta_0}{6!} \int d^D x g^3(x, x) + O(1), \quad (15)$$

where the trace is only in the functional space, and the

gap equation becomes

$$g^{-1}(x, y) = i \left[\square + m_0^2 + \frac{\lambda_0}{6} \left(\frac{\phi^2}{N} + g(x, x) \right) + \frac{\eta_0}{5!} \left(\frac{\phi^2}{N} + g(x, x) \right)^2 \right] \delta^D(x - y) + O\left(\frac{1}{N}\right). \quad (16)$$

It is important to point out that this calculation was done by Townsend [9]. We interested to generalize these results assuming that the system is in equilibrium with a thermal bath a temperature $T = \beta^{-1}$. Since we are studying the equilibrium situation it is convenience to use the Matsubara formalism. Consequently it is convenient to continue all momenta to Euclidean values

($p_0 = ip_4$) and take the following Ansatz for $g(x, y)$,

$$g(x, y) = \int \frac{d^D p}{(2\pi)^D} \frac{\exp^{i(x-y)p}}{p^2 + M^2(\phi)}. \quad (17)$$

Substituting eq.(17) in eq.(16) we get the expression for

the gap equation:

$$M^2(\phi) = m_0^2 + \frac{\lambda_0}{6} \left(\frac{\phi^2}{N} + F(\phi) \right) + \frac{\eta_0}{5!} \left(\frac{\phi^2}{N} + F(\phi) \right)^2, \quad (18)$$

where $F(\phi)$ is given by

$$F(\phi) = \int \frac{d^D p}{(2\pi)^D} \frac{1}{p^2 + M^2(\phi)}, \quad (19)$$

and the effective potential in the D-dimensional Euclidean space can be expressed as

$$V(\phi) = V_0(\phi) + \frac{N}{2} \int \frac{d^D p}{(2\pi)^D} \ln [p^2 + M^2(\phi)] - \frac{N}{4!} (\lambda_0 + \frac{\eta_0 \phi^2}{10N}) F(\phi)^2 - \frac{2N \eta_0 F(\phi)^3}{6!}, \quad (20)$$

where $V_0(\phi)$ is the classical potential. In the next section using the Matsubara formalism we present the effective potential of the model at finite temperature.

3 The effective potential for φ^6 theory at finite temperature

Let us suppose that our system is in equilibrium with a thermal bath. To study the temperature effects in quantum field theory we will use the imaginary time Green function approach [10]. In this formalism the Euclidean-time τ is restricted to the interval $0 \leq \tau \leq \beta$, and the bosonic field satisfies periodic boundary conditions in Euclidean-time. This is equivalence to replace

the continuous four momenta k_4 by discrete ω_n and the integration by a summation ($\beta = \frac{1}{T}$):

$$k_4 \rightarrow \omega_n = \frac{2\pi n}{\beta}, \quad n = 0, \pm 1, \pm 2, \dots$$

$$\int \frac{d^D k}{(2\pi)^D} \rightarrow \sum_n \frac{1}{\beta} \int \frac{d^{D-1} k}{(2\pi)^{D-1}}. \quad (21)$$

It is important to stress that all the Feynman rule are the same as the temperature case, except, as we stressed that momentum space integrals over the zeroth component is replace by summ over discret sums. The effective potential at finite temperature can be write as:

$$V_\beta(\phi) = V_0(\phi) + \frac{N}{2\beta} \sum_n \int \frac{d^{D-1} p}{(2\pi)^{D-1}} \ln [\omega_n^2 + p^2 + M_\beta^2(\phi)] - \frac{N}{4!} (\lambda_0 + \frac{\eta_0 \phi^2}{10N}) F_\beta(\phi)^2 - \frac{2N \eta F_\beta(\phi)^3}{6!}. \quad (22)$$

where $F_\beta(\phi)$ is a finite temperature generalization of $F(\phi)$, where,

$$F_\beta(\phi) = \frac{1}{\beta} \sum_{n=-\infty}^{\infty} \int \frac{d^{D-1} p}{(2\pi)^{D-1}} \frac{1}{\omega_n^2 + p^2 + M_\beta^2(\phi)}. \quad (23)$$

The gap equation for this theory at finite temperature is given by,

$$M_\beta^2(\phi) = m_0^2 + \frac{\lambda_0}{6} \left(\frac{\phi^2}{N} + F_\beta(\phi) \right) + \frac{\eta_0}{5!} \left(\frac{\phi^2}{N} + F_\beta(\phi) \right)^2. \quad (24)$$

In order to regularize $F_\beta(\phi)$ given by eq.(23), we use a mixing between dimensional regularization and analytic regularization. For this purpose we define the expression $I_\beta(D, s, m)$ as :

$$I_\beta(D, s, m) = \frac{1}{\beta} \sum_{n=-\infty}^{\infty} \int \frac{d^{D-1} k}{(2\pi)^{D-1}} \frac{1}{(\omega_n^2 + k^2 + m^2)^s}. \quad (25)$$

Using the analytic extension of the inhomogeneous Epstein zeta function it is possible to obtain the analytic extension of $I_\beta(D, s, m)$;

$$I_\beta(D, s, m) = \frac{m^{D-2s}}{(2\pi^{1/2})^D \Gamma(s)} \left[\Gamma(s - \frac{D}{2}) + 4 \sum_{n=1}^{\infty} \left(\frac{2}{mn\beta} \right)^{D/2-s} K_{D/2-s}(mn\beta) \right] \quad (26)$$

where $K_\mu(z)$ is the modified Bessel function of the third kind. Fortunately for $D = 3$ the analytic extension of the function $I_\beta(D, s = 1, m = M_\beta(\phi)) = F_\beta(\phi)$ is finite

and can be expressed in a closed form [5] (note that in $D = 3$ we have no pole, at least in this approximation), and in particular as

$$F_\beta(\phi) = I_\beta(3, 1, M_\beta(\phi)) = -\frac{M_\beta(\phi)}{4\pi} \left(1 + \frac{2 \ln(1 - e^{-M_\beta(\phi)\beta})}{M_\beta(\phi)\beta} \right). \quad (27)$$

This result is not a peculiarity of this method of regularization, because this happens also in dimensional regularization at zero temperature, that is, in odd dimensions integrals which are divergent by naive power counting may be regulated to finite value with no poles occurring, for example in $D = 3$. In order to regularize the second term of eq.(22), we use the following method: We define,

$$LF_\beta(\phi) = \frac{1}{\beta} \sum_{n=1}^{\infty} \int \frac{d^{D-1}p}{(2\pi)^{D-1}} \ln [\omega_n + p^2 + M_\beta^2(\phi)] \quad (28)$$

then,

$$\frac{\partial LF_\beta(\phi)}{\partial M_\beta} = (2M_\beta) \frac{1}{\beta} \sum_{n=1}^{\infty} \int \frac{d^{D-1}p}{(2\pi)^{D-1}} \frac{1}{\omega_n + p^2 + M_\beta^2(\phi)} \quad (29)$$

and from the equation (23), we have that,

$$\frac{\partial LF_\beta(\phi)}{\partial M_\beta} = (2M_\beta) F_\beta(\phi), \quad (30)$$

in this way the effective potential could be regularized. For $D = 3$, $F_\beta(\phi)$ is finite and is given by eq. (31)[5] and integrating the eq.(30), we obtain:

$$LF_\beta(\phi)_R = -\frac{M_\beta(\phi)^3}{6\pi} - \frac{M_\beta(\phi) Li_2(e^{-M_\beta(\phi)\beta})}{\pi \beta^2} - \frac{Li_3(e^{-M_\beta(\phi)\beta})}{\pi \beta^3}. \quad (31)$$

The definition of general polylogarithm function $Li_n(z)$ can be found in ref. [11].

The daisy and super daisy resummed effective potential at finite temperature for $D = 3$ is given by:

$$V_\beta(\phi) = V_0(\phi) + \frac{N}{2} LF_\beta(\phi)_R - \frac{N}{4!} (\lambda_0 + \frac{\eta_0 \phi^2}{10N}) (F_\beta(\phi)_R)^2 - \frac{2N \eta (F_\beta(\phi)_R)^3}{6!}. \quad (32)$$

and the gap equation (see eq.(24)):

$$M_\beta^2(\phi) = m_0^2 + \frac{\lambda_0}{6} \left(\frac{\phi^2}{N} - \frac{M_\beta(\phi)}{4\pi} \left[1 + \frac{2 \ln(1 - e^{-M_\beta(\phi)\beta})}{M_\beta(\phi)\beta} \right] \right) + \frac{\eta_0}{5!} \left(\frac{\phi^2}{N} - \frac{M_\beta(\phi)}{4\pi} \left[1 + \frac{2 \ln(1 - e^{-M_\beta(\phi)\beta})}{M_\beta(\phi)\beta} \right] \right)^2. \quad (33)$$

The numerical result for the effective potential for $D = 3$ at three different temperature is reported in

Fig.(2).

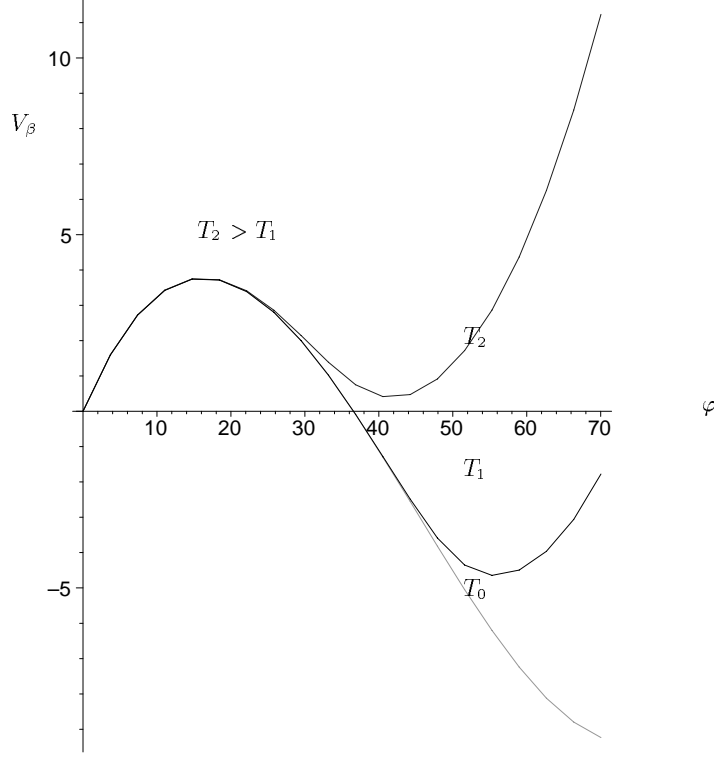


Figure 2.: The daisy and superdaisy resummed effective potential at three different temperature $T_2 > T_1 > T_0$ for $D = 3$.

From the figure below, we see that, there exist a first order transition for the parameter of the theory in the case when the effective potential at tree level, presents symmetry breaking at zero temperature (denoted for curve label by T_0).

For the case $D = 4$, and for $\eta = 0$, where the theory is just renormalizable, the effective potential can be obtain in the same way, that is:

$$V_\beta(\phi) = V_0(\phi) + \frac{N}{2}(LF_\beta(\phi)_R)_{D=4} - \frac{\lambda_0 N}{4!}(F_\beta(\phi)_R)_{D=4}^2, \quad (34)$$

where V_0 is the classical potential,

$$V_0(\phi) = \frac{1}{2}m^2\phi^2 + \frac{\lambda}{4!N}\phi^4, \quad (35)$$

and $(F_\beta(\phi)_R)_{D=4}$ is given by:

$$(F_\beta(\phi)_R)_{D=4} = \frac{\lambda M_\beta^2(\phi)}{2\pi^2} \int_1^\infty \frac{(p^2 - 1)^{\frac{1}{2}}}{e^{M_\beta(\phi)\beta p} - 1} dp \quad (36)$$

and in the limit of high temperature, we could write the above equation as:

$$(F_\beta(\phi)_R)_{D=4} = \frac{1}{12\beta^2} - \frac{M_\beta(\phi)}{4\pi\beta} - \frac{M_\beta^2(\phi)}{8\pi^2} \ln(M_\beta(\phi)\beta); \quad (37)$$

and,

$$(LF_\beta(\phi)_R)_{D=4} = \frac{M_\beta(\phi)^2}{12\beta^2} - \frac{M_\beta(\phi)^3}{6\pi\beta} - \frac{M_\beta(\phi)^4 \ln(M_\beta(\phi)\beta)}{16\pi^2} + \frac{M_\beta(\phi)^4}{64\pi^2}, \quad (38)$$

and the gap equation for $D = 4$ is given by:

$$M_\beta^2(\phi) = \tilde{m}^2(\phi) + \frac{\lambda}{6} \left(\frac{1}{12\beta^2} - \frac{M_\beta(\phi)}{4\pi\beta} - \frac{M_\beta^2(\phi)}{8\pi^2} \ln(M_\beta(\phi)\beta) \right), \quad (39)$$

where $\tilde{m}^2(\phi) = m^2 + \frac{\lambda}{6} \frac{\phi^2}{N}$, and m^2 , λ are the renormalized mass and coupling constant a temperature zero respectively. We note that, from the gap equation in eq.(39), we find that for the coupling constant $\lambda \ll 1$, the condition $M_\beta^2(\phi)/T^2 \ll 1$ is consistent with $\tilde{m}^2(\phi)/T^2$, which is exactly the required condition for the high temperature expansion [10].

4 Conclusions

In this paper we have performed an analysis of the daisy and super daisy effective potential for the theory φ^4 and φ^6 in D -dimensional Euclidean space at finite temperature. The form of effective potential have been found explicitly using resummation method in the leading order $1/N$ approximation (Hartree-Fock approximation). We have seen how dimensional regularization and analytic regularization can be used to compute the effective potential at finite temperature. In odd dimensional theory when power counting indicates that the diverges should occur, dimensional regularization and analytic does not give rise to a pole.

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