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BEHAVIOUR OF COUPLING CONSTANTS AT HIGH TEMPERATURE IN SUPERSYMMETRIC THEORIES *

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ABSTRACT

An analysis is presented of the temperature dependence of the coupling constants using the improved one-loop approximation in the Wess-Zumino model and the supersymmetric O(N) model. It is found that all the coupling constants, both bosonic (ϕ^4 type) and Yukawa, approach constant nonzero values as $T \to \infty$. The asymptotic values of the bosonic couplings are slightly smaller than the corresponding zero-temperature values, and those of the Yukawa couplings are the same as the zero-temperature values.

MIRAMARE - TRIESTE
September 1986

* To be submitted for publication.

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Recently the behaviour of the coupling constant at high temperature has been investigated within the framework of the $\lambda\phi^4$ theory[1)-4)], and it is found to possess the surprising property that it approaches a constant non-zero value as $T \to \infty$, rather than decreasing to zero. This conclusion is reached independently by two different methods, the improved one-loop approximation[1),3),4)], and the calculation based on the Dyson-Schwinger equation[2)].

The improved one-loop approximation was first employed by Dolan and Jackiw[4)] to compute the temperature-dependent mass to the next-to-leading order in T, i.e., the linear T term. It is well-known that simple one-loop approximation does not give reliable results for the temperature-dependent coupling constant[5)]. This is because the temperature dependence of the coupling constant is only linear in T. These linear T contributions are known to be strongly modified by higher-loop effects[1)].

In the improved one-loop approximation[1),3),4)], the higher-loop contributions are effectively taken into account by replacing each of the internal bare propagators in the one-loop approximation by a dressed propagator. The mass of the dressed propagator, which is temperature-dependent, is determined by a self-consistency relation. Upon iteration, the self-consistency relation can be easily shown to yield a summation of all higher-loop diagrams with no overlapping loops[1),3),4)].

For the case of one-component $\lambda\phi^4$ theory, the temperature-dependent coupling constant is given by[1)-4)]

$$\lambda(T) = \lambda - \frac{3\lambda^2}{16\pi}\left(\frac{T}{\mu_T} - \frac{1}{\pi}\ln\frac{T}{\mu_T}\right), \tag{1}$$

where μ_T is the temperature-dependent mass used in the dressed propagators, and is given by

$$\mu_T^2 = \mu^2 + \frac{\lambda}{24}T^2. \tag{2}$$

For very high T, it is easily seen that[1)-4)]

$$\lambda(T) \longrightarrow \lambda - \frac{3\lambda^2}{16\pi}\left[\left(\frac{24}{\lambda}\right)^{1/2} + \frac{1}{2\pi}\ln\frac{\lambda}{24}\right]. \tag{3}$$

One point is worth noting here. In the improved one-loop approximation, we are summing an infinite series of diagrams. The renormalization prescription at $T \neq 0$ may be different from that at $T = 0$. In fact, in order to remove the divergences in a self-consistent way, we are forced into a temperature-dependent renormalization[4]. If we write the counter terms in the $\lambda\phi^4$ Lagrangian as

$$\mathcal{L}_C = -\frac{1}{2} B\phi^2 - \frac{1}{4!} C\phi^4, \tag{4}$$

then the renormalization constant B and C are respectively given by

$$B = -\frac{1}{2}\lambda \int \frac{d^4k}{(2\pi)^4} \frac{-i}{k^2+\mu_T^2+i\epsilon}, \tag{5}$$

$$C = \frac{3}{2} i\lambda^2 \int \frac{d^4k}{(2\pi)^4} \left(\frac{-i}{k^2+\mu_T^2+i\epsilon}\right)^2. \tag{6}$$

The metric used in this paper is (-1, 1, 1, 1).

In this paper, we shall extend the analysis to deal with theories with global supersymmetry. Our purpose is two-fold. First of all, we want to find out whether the presence of fermions in a specific way (as demanded by supersymmetry) would affect the behaviour of the 4-point bosonic coupling constants at high T. Secondly, in a supersymmetric theory, the Yukawa coupling constant is related to the 4-point bosonic coupling constant at $T = 0$. For example, for the Wess-Zumino model [7], only one coupling constant is required to describe both the 4-point bosonic coupling as well as the Yukawa coupling. How would the Yukawa coupling constant in a supersymmetric theory behave at high T?

To begin with, let us consider the simplest supersymmetric theory, the Wess-Zumino model [7]. We choose, without loss of generality, the following superpotential:

$$f_{WZ} = -\lambda S + \frac{1}{6} g S^3. \tag{7}$$

Hereafter, we shall use the notations A, B and ψ to denote the scalar, the pseudoscalar and the spinor components of the superfield S. The Wess-Zumino model is known to undergo a second-order phase transition [8] at $T_c = (8\lambda/g)^{1/2}$. For $T > T_c$, the system is characterized by

$$\langle A \rangle_T = 0, \tag{8}$$

where $\langle\ \rangle_T$ denotes ordinary thermal average. The Lagrangian for the Wess-Zumino model, in terms of the component fields, is

$$\mathcal{L} = -\frac{1}{2}\partial_\mu A \partial^\mu A - \frac{1}{2}\partial_\mu B \partial^\mu B - \frac{1}{2}\bar{\psi}\not{\partial}\psi - \frac{1}{2}\mu_A^2 A^2 - \frac{1}{2}\mu_B^2 B^2$$
$$- \frac{g^2}{16}(A^2+B^2)^2 - \frac{g}{2\sqrt{2}}\bar{\psi}(A+iB\gamma_5)\psi, \tag{9}$$

where the tree-level masses are

$$\mu_{A,B}^2 = \mp \lambda g, \qquad \mu_\psi = 0. \tag{10}$$

We observe that the same quantity g describes both the 4-point bosonic as well as the Yukawa couplings at the tree-level. We further notice that μ_A^2 is negative. Thus the coupling constants as calculated in the simple one-loop approximation would yield complex quantities. In order to apply the improved one-loop approximation, we have to compute first the temperature-dependent masses. This is easily performed, and we obtain the following temperature-dependent masses for $T > T_c$:

$$\tilde{\mu}_{A,B}^2 = \mu_{A,B}^2 + \frac{1}{8} g^2 T^2, \qquad \tilde{\mu}_\psi = 0. \tag{11}$$

Throughout this paper, we shall use the notation that a symbol with a tilde denotes the corresponding temperature-dependent quantity at high T. We observe that ψ is massless at all $T > T_c$. It is the massless Goldstone fermion state [9].

We now proceed to compute the temperature-dependent coupling constants. We consider first the A^4 self-coupling constant. The diagrams contributed to the A^4 coupling constant are as depicted in Figs. 1(a) and 1(b). The boson loop in Fig. 1(a) can be A or B. The evaluation of the boson loop contributions is straightforward, and is similar to the case of the N-component $\lambda\phi^4$ theory [1),3)]. But the fermion loop diagram, Fig. 1(b), is infra-red divergent because the fermion is massless. A proper treatment of the temperature-dependent infra-red divergence problem will entail a detailed analysis of the processes of emission and absorption of particles from the heat bath[10]. In our present case, the fermion loop diagram of Fig. 1(b) is proportional to

$$-\frac{g^4}{64}\int_k \mathrm{Tr}\left[\frac{-i\not{k}}{k^2}\right]_f \left[\frac{-i\not{k}}{k^2}\right]_f \left[\frac{-i\not{k}}{k^2}\right]_f \left[\frac{-i\not{k}}{k^2}\right]_f = \frac{g^4}{16}\int_k \left[\frac{-i}{k^2}\right]_f \left[\frac{-i}{k^2}\right]_f . \tag{12}$$

This contribution is clearly infra-red divergent. This divergence is, however, cancelled by the process in which the A-field absorbs and re-emits soft massless fermions from the heat bath. This absorption and re-emission process has the same structure as Eq. (12), but has the opposite sign, thus cancelling the fermion-loop contribution from Fig. 1(b). Upon removal of the infra-red divergence, we arrive at the following expression for the A^4 coupling constant at high T,

$$\tilde{g}^2(A^4) = g^2\left\{1 - \frac{9g^2}{32\pi}F_1(T/\tilde{\mu}_A) - \frac{g^2}{32\pi}F_1(T/\tilde{\mu}_B)\right\}, \tag{13}$$

where $\tilde{\mu}_A$ and $\tilde{\mu}_B$ are as given in (11), and the function $F_1(x)$ is defined by

$$F_1(x) = x - \frac{1}{\pi}\ln x . \tag{14}$$

As $T \to \infty$, both $T/\tilde{\mu}_A$ and $T/\tilde{\mu}_B$ are large, and become equal in the limit. Thus,

$$\tilde{g}^2(A^4) \xrightarrow[T\to\infty]{} g^2\left(1 - \frac{g}{4\sqrt{3}\,\pi} - \frac{g^2}{16\pi^2}\ln g\right). \tag{15}$$

The A^4 coupling constant, therefore, approaches a nonzero value as $T \to \infty$, with the asymptotic value being slightly below the tree-level value. This is in agreement with the behaviour of the coupling constant obtained previously[1)-4)] in the N-component $\lambda\phi^4$ theory. This result is not surprising at all. This is because only the boson loops contribute to the A^4 coupling constant. We thus expect the behaviour to be similar. The fermion loop contribution, being infra-red divergent, is removed by hand. The B^4-coupling and the A^2B^2 coupling can be similarly calculated. They are given by expressions similar to, but slightly different from, that of Eq. (13). They, however, approach the same asymptotic value, Eq.(15), as $T \to \infty$.

The Yukawa coupling constants are calculated from diagrams of Fig 1(c). For the coupling of A to ψ, we have two diagrams contributing. The internal boson in Fig. 1(c) may be an A or a B. The $\bar{\psi}A\psi$ coupling constant is therefore given by

$$\begin{aligned}\tilde{g}(\bar{\psi}A\psi) &= g - \tfrac{1}{2}ig^2\bar{\int_k}\left[\frac{-i}{k^2}\right]_f\left\{\left[\frac{-i}{k^2+\tilde{\mu}_A^2}\right]_b - \left[\frac{-i}{k^2+\tilde{\mu}_B^2}\right]_b\right\},\\ &= g - \tfrac{1}{2}g^2\left(\frac{1}{\tilde{\mu}_A^2} - \frac{1}{\tilde{\mu}_B^2}\right)\bar{\int_k}\left[\frac{-i}{k^2}\right]_f\\ &\quad + \tfrac{1}{2}g^2\left\{\frac{1}{\tilde{\mu}_A^2}\bar{\int_k}\left[\frac{-i}{k^2+\tilde{\mu}_A^2}\right]_b - \frac{1}{\tilde{\mu}_B^2}\bar{\int_k}\left[\frac{-i}{k^2+\tilde{\mu}_B^2}\right]_b\right\}.\end{aligned} \tag{16}$$

The integration $\int_k$ has the usual meaning of the momentum integration at $T \neq 0$, and the bar above it signifies that only the finite-temperature part of the momentum integration is to be retained. The subscripts f and b on the propagators denote respectively fermion and boson propagators. The fermion propagator is written without the spin projection. The integrations over the boson propagators in (16)are straightforward. For the integration over the fermion propagator, we use the formula

$$\bar{\int_k}\left[\frac{-i}{k^2+m^2}\right]_f = -\frac{T^2}{24} + \frac{m^2}{16\pi^2}\ln\frac{T^2}{m^2} + \frac{1+2\ln\pi-2\gamma}{16\pi^2}m^2 + \cdots . \tag{17}$$

Putting all the factors together, we arrive at the following expression for the $\bar{\psi}A\psi$ coupling constant

$$\tilde{g}(\bar{\psi}A\psi) = g\left\{1 + \frac{g^2}{16}F_2(T/\tilde{\mu}_A) - \frac{g^2}{16}F_2(T/\tilde{\mu}_B)\right\}, \tag{18}$$

where the function $F_2(x)$ is defined by

$$F_2(x) = x^2 - \frac{2}{\pi}x + \frac{1}{2\pi^2}\ln x ; \tag{19}$$

As $T \to \infty$, since both $\tilde{\mu}_A^2$, $\tilde{\mu}_B^2 \to \frac{1}{8}g^2T^2$, and the temperature dependence in (18) is of order $1/T^2$. Therefore,

$$\tilde{g}(\bar{\psi}A\psi) \longrightarrow g , \qquad \text{as } T \to \infty . \tag{20}$$

The coupling of B to ψ is similarly calculated. Denoting the coupling constant at the $\bar{\psi} i\gamma_5 B\psi$ vertex by $\tilde{g}(\bar{\psi} i\gamma_5 B\psi)$, we have

$$\tilde{g}(\bar{\psi} i\gamma_5 B\psi) = \tilde{g}(\bar{\psi} A\psi) \quad \text{with} \quad \tilde{\mu}_A \longleftrightarrow \tilde{\mu}_B \tag{21}$$

$$\longrightarrow g \quad \text{as} \quad T \longrightarrow \infty . \tag{22}$$

This result is significant. Although the Yukawa coupling constant also approaches nonzero value as $T \to \infty$, the asymptotic value is nevertheless equal to the tree-level value. This implies that the 4-point bosonic coupling constant and the Yukawa coupling constant do behave differently at high T.

Next we consider the supersymmetric $O(N)$ model[8),11)] which is obtained by coupling the Wess-Zumino model to an N-component chiral superfield. The superpotential for this model is given by

$$f_{O(N)} = -\lambda S_0 + \frac{1}{6} g S_0^3 + \frac{1}{2} h S_0 (S_1^2 + S_2^2 + \cdots + S_N^2) . \tag{23}$$

There are two coupling constants: g to describe the self-coupling of S_0, and h to describe the coupling of S_0 to S_i ($i = 1, 2, \ldots N$). This model is interesting on its own. It exhibits the phenomena of symmetry anti-restoration[12)], for it undergoes a first-order phase transition followed by a second-order phase transition[8),11)]. For $T > T_c = (4\lambda/h)^{1/2}$, the system is characterized by[8),11)]

$$\langle A_0 \rangle_T = \langle A_1 \rangle_T = \langle A_2 \rangle_T = \cdots = \langle A_N \rangle = 0 . \tag{24}$$

The Lagrangian of the system contains an interaction part given by

$$\begin{aligned}\mathcal{L}_{int} = &- \frac{g}{2\sqrt{2}} \bar{\psi}_0 (A_0 + i\gamma_5 B_0)\psi_0 - \frac{h}{2\sqrt{2}} \sum_i \bar{\psi}_i (A_0 + i\gamma_5 B_0)\psi_i - \frac{h}{\sqrt{2}} \sum_i \bar{\psi}_i (A_i + i\gamma_5 B_i)\psi_i \\ &- \frac{g}{16}(A_0^2 + B_0^2)^2 - \frac{h^2}{16}\Big[\sum_i (A_i^2 - B_i^2)\Big]^2 - \frac{h^2}{4}\Big[\sum_i A_i B_i\Big]^2 - \frac{1}{2} gh A_0 B_0 \sum_i A_i B_i \\ &- \frac{1}{8} h(2h+g) \sum_i (A_0^2 A_i^2 + B_0^2 B_i^2) - \frac{1}{8} h(2h-g) \sum_i (A_0^2 B_i^2 + B_0^2 A_i^2) .\end{aligned} \tag{25}$$

The tree-level masses are given by

$$\mu_{A_0, B_0}^2 = \mp \lambda g , \qquad \mu_{A_i, B_i}^2 = \mp \lambda h , \qquad \mu_{\psi_0} = \mu_{\psi_i} = 0 . \tag{26}$$

The index i in (25) and (26) is understood to run from 1 to N. The computation of the temperature-dependent masses is again straightforwardly performed, and we obtain

$$\left.\begin{aligned}\tilde{\mu}_{A_0, B_0}^2 &= \mu_{A_0, B_0}^2 + \frac{1}{8}(g^2 + N h^2) T^2 , \\ \tilde{\mu}_{A_i, B_i}^2 &= \mu_{A_i, B_i}^2 + \frac{1}{4} h^2 T^2 , \\ \tilde{\mu}_{\psi_0} &= \tilde{\mu}_{\psi_i} = 0 .\end{aligned}\right\} \tag{27}$$

All the fermions are massless for $T > T_c$.

The computation of the temperature-dependent coupling constants in the supersymmetric $O(N)$ model proceeds in essentially the same way as for the case of the Wess-Zumino model. Here, the number of diagrams contributing is much larger. However, one feature is common in both cases, i.e., the fermions are massless. The contributions to the bosonic coupling constants from diagrams involving fermion loops will therefore be infra-red divergent.
Here again, such infra-red divergences are removed by processess of absorption and re-emission of soft massless fermions from the heat bath, similar to the previous case. Instead of presenting all the coupling constants calculated, we choose to present here only a representative few, namely those at the following couplings: A_0^4, A_i^4, $\bar{\psi}_0 A_0 \psi_0$ and $\bar{\psi}_i A_0 \psi_i$. At the tree-level, the A_0^4 and the $\bar{\psi}_0 A_0 \psi_0$ couplings represent respectively the 4-point bosonic coupling and the Yukawa coupling described by the same coupling constant g; and the A_i^4 and $\bar{\psi}_i A_0 \psi_i$ couplings represent those described by h. The coupling constants for these couplings, as calculated from the improved one-loop approximation, are

$$\begin{aligned}\tilde{g}^2(A_0^4) = g^2 \Big\{ 1 &- \frac{9g^2}{32\pi} F_1(T/\tilde{\mu}_{A_0}) - \frac{g^2}{32\pi} F_1(T/\tilde{\mu}_{B_0}) \\ &- \frac{N h^2}{32\pi g^2} \Big[(2h+g)^2 F_1(T/\tilde{\mu}_{A_i}) + (2h-g)^2 F_1(T/\tilde{\mu}_{B_i}) \Big] \Big\} ,\end{aligned} \tag{28}$$

$$\tilde{h}^2(A_i{}^4) = h^2 \left\{ 1 - \frac{(2h+g)^2}{32\pi} F_1(T/\tilde{\mu}_{A_0}) - \frac{(2h-g)^2}{32\pi} F_1(T/\tilde{\mu}_{B_0}) \right.$$

$$\left. - \frac{(N+8)h^2}{32\pi} F_1(T/\tilde{\mu}_{A_i}) - \frac{Nh^2}{32\pi} F_1(T/\tilde{\mu}_{B_i}) \right\} , \tag{29}$$

$$\tilde{g}(\bar{\psi}_0 A_0 \psi_0) = g \left\{ 1 + \frac{g^2}{16} \left[F_2(T/\tilde{\mu}_{A_0}) - F_2(T/\tilde{\mu}_{B_0}) \right] \right.$$

$$\left. + \frac{Nh^3}{16g} \left[F_2(T/\tilde{\mu}_{A_i}) - F_2(T/\tilde{\mu}_{B_i}) \right] \right\} , \tag{30}$$

$$\tilde{h}(\bar{\psi}_i A_0 \psi_i) = h \left\{ 1 + \frac{h^2}{16} \left[F_2(T/\tilde{\mu}_{A_0}) - F_2(T/\tilde{\mu}_{B_0}) \right] \right.$$

$$\left. + \frac{gh}{16} \left[F_2(T/\tilde{\mu}_{A_i}) - F_2(T/\tilde{\mu}_{B_i}) \right] \right\} . \tag{31}$$

The functions $F_1(x)$ and $F_2(x)$ are as defined by Eqs.(14) and (19), and the masses are the temperature-dependent masses given by Eq.(27). Noting that all the boson masses increase linearly with T, we obtain the following nonzero asymptotic values for the coupling constants as $T \to \infty$:

$$\tilde{g}^2(A_0{}^4) \longrightarrow g^2 \left\{ 1 - \frac{N(g^2+4h^2)}{8\pi g^2} \left(h + \frac{h^2}{2\pi} \ln h \right) \right.$$

$$\left. - \frac{5g}{4\sqrt{2}\,\pi} \left(1 + \frac{Nh^2}{g^2} \right)^{-1/2} - \frac{5g^2}{32\pi^2} \ln (g^2 + Nh^2) \right\} , \tag{32}$$

$$\tilde{h}^2(A_i{}^4) \longrightarrow h^2 \left\{ 1 - \frac{N+4}{8\pi} \left(h + \frac{h^2}{2\pi} \ln h \right) - \frac{1}{4\sqrt{2}\,\pi} (g^2+4h^2)(g^2+Nh^2)^{-1/2} \right.$$

$$\left. - \frac{1}{32\pi^2} (g^2+4h^2) \ln (g^2+Nh^2) \right\} . \tag{33}$$

$$\tilde{g}(\bar{\psi}_0 A_0 \psi_0) \longrightarrow g , \tag{34}$$

$$\tilde{h}(\bar{\psi}_i A_0 \psi_i) \longrightarrow h . \tag{35}$$

The results, as given by Eqs.(32)-(35), are very similar to those obtained for the Wess-Zumino model. All coupling constants approach nonzero values as $T \to \infty$. The asymptotic values for the 4-point bosonic couplings are slightly smaller than the corresponding tree-level values, whereas those for the Yukawa couplings are the same as the corresponding tree-level values.

In summary, we have presented the analysis of the temperature dependence of the coupling constants using the improved one-loop approximation in two supersymmetric theories, the Wess-Zumino model and the supersymmetric O(N) model. In a supersymmetric theory, the 4-point bosonic couplings and the Yukawa couplings are generally related at $T = 0$, with the same constant describing both types of couplings. Our analysis shows that all coupling constants, the 4-point bosonic as well as the Yukawa, approach nonzero values as $T \to \infty$. But the asymptotic behaviours are different for the two types of couplings. The asymptotic values for the 4-point bosonic couplings are slightly smaller than the corresponding zero-temperature values, whereas those for the Yukawa couplings are the same as the zero-temperature values. The results for the 4-point bosonic couplings are similar to that obtained previously[1)-4)] for the N-component $\lambda\phi^4$ theory. It is somewhat expected because the temperature dependence of the bosonic coupling constants receives contributions from diagrams involving only boson loops. The contributions from diagrams with fermion loops are infra-red divergent, but are removed through considerations of processes of absorption and emission of particles from the heat bath. The Yukawa couplings do not involve any infra-red divergent contributions. The result that they all approach to the respective zero-temperature values is therefore significant.

ACKNOWLEDGMENTS

The author would like to thank Professor Abdus Salam, the International Atomic Energy Agency and UNESCO for hospitality at the International Centre for Theoretical Physics, Trieste.

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FIGURE CAPTION

Fig. 1 Diagrams contributing to the temperature dependence of the coupling constants. Fermions are denoted by solid lines, and bosons by broken lines. The internal propagators are understood to be dressed propagators. Bosonic couplings receive contributions from (a) boson loop diagram, and (b) fermion loop diagram. Yukawa couplings receive contributions only from (c).

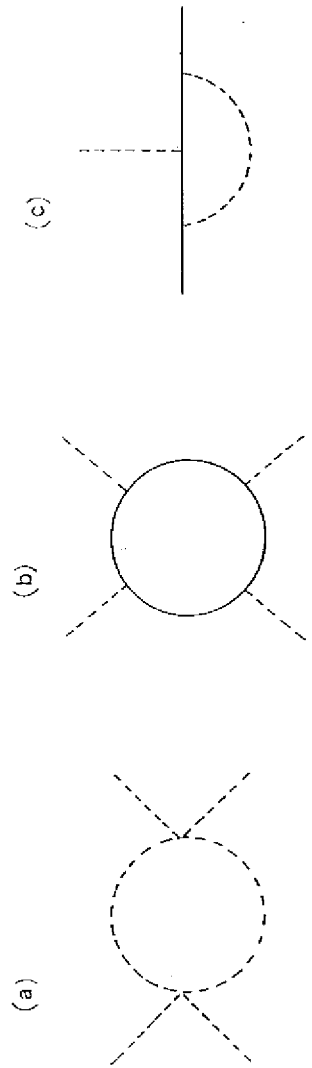

Fig. 1