# Five-Loop Critical Temperature Shift in Weakly Interacting Homogeneous Bose-Einstein Condensate

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Using variational perturbation theory, we calculate the shift in the critical temperature up to five loops to lowest order in the scattering length $\Delta T_c/T_c \approx 0.93 \pm 0.13\ an^{1/3}$.

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The effect of a small repulsive interaction upon the critical temperature $T_c$ of a Bose-Einstein condensate (BEC) has been a matter of controversy for many years, and the various results are converging only very slowly towards a common answer. Here we want to contribute the hopefully most accurate result so far.

The interacting Bose-Einstein condensate is described by the euclidean action

$$\mathcal{A}_E = \int_0^\beta d\tau \int d^3x \left\{ \psi^*(\mathbf{x},\tau) \left( \partial_\tau - \frac{1}{2M}\nabla^2 \right) \psi(\mathbf{x},\tau) - \mu\psi^*(\mathbf{x},\tau)\psi(\mathbf{x},\tau) + \frac{2\pi a}{M} \left[\psi(\mathbf{x},\tau)\psi^*(\mathbf{x},\tau)\right]^2 \right\} , \quad (1)$$

where $M$ is the mass of the bosons, $\beta$ the inverse temperature in natural units with $\hbar = k_B = 1$, $a$ is the $s$-wave scattering length, and $\mu$ the chemical potential. The free system has a transition temperature

$$T_c^{(0)} = \frac{2\pi}{M} \left[ \frac{n}{\zeta(3/2)} \right]^{\frac{2}{3}} , \quad (2)$$

where $n$ is the particle density. A small relative shift of $T_c$ with respect to $T_c^{(0)}$ can be calculated from the general formula

$$\frac{\Delta T_c}{T_c^{(0)}} = -\frac{2}{3}\frac{\Delta n}{n^{(0)}}, \quad (3)$$

where $n^{(0)}$ is the particle density in the free condensate and $\Delta n$ its change at $T_c$ caused by the small interaction. For small $a$, this behaves like [1, 2]

$$\frac{\Delta T_c}{T_c^{(0)}} = c_1 a n^{1/3} + [c_2' \ln(an^{1/3}) + c_2] a^2 n^{2/3} + \mathcal{O}(a^3 n). \quad (4)$$

where $c_2' = -64\pi\zeta(1/2)/3\zeta(3/2)^{5/3} \simeq 19.7518$ can be calculated perturbatively, whereas $c_1$ and $c_2$ require nonperturbative techniques since infrared divergences at $T_c$ make them basically strong-coupling results. The standard technique to reach this regime is based on a resummation of perturbation expansions using the renormalization group [3, 4], first applied in this context by Ref. [6]. Recently, however, it has been shown by calculating the best known critical exponent $\alpha$ of superfluid helium from Satellite experiments [7] that the accuracy of strong-coupling results can be surpassed by much simpler variational perturbation theory [4, 8, 9].

Up to now, $c_2$ has been inferred only from Monte Carlo data to be $c_2 \approx 75.7 \pm 0.4$. In order to find the leading coefficient $c_1$, one may take advantage of an important simplification due to the fact that $\Delta n$ can be calculated from the classical limit of the field theory, which is governed by the three-dimensional action

$$\mathcal{A}_{3d} = \beta \int d^3x \left\{ \psi_0^*(\mathbf{x}) \left( -\frac{1}{2M}\nabla^2 - \mu \right) \psi_0(\mathbf{x}) + \frac{2\pi a}{M} \left[\psi_0^*(\mathbf{x})\psi_0(\mathbf{x})\right]^2 \right\} . \quad (5)$$

This is a special case $N = 2$ of the more general O($N$)-invariant $\phi^4$ field theory

$$\mathcal{A}_\phi = \int d^3x \left[ \frac{1}{2}|\nabla\phi|^2 + \frac{1}{2}m^2\phi^2 + \frac{u}{4!}(\phi^2)^2 \right] , \quad (6)$$

where the $N$-component field $\phi = (\phi_1, \phi_2, \ldots, \phi_N)$ is related to the original field $\psi$ for $N = 2$ by $\psi(\mathbf{x}) = \sqrt{MT}[\phi_1(\mathbf{x}) + i\phi_2(\mathbf{x})]$. The square mass is $m^2 = -2M\mu$, and the quartic coupling is $u = 48\pi aMT$. Using this relation, the shift of the critical temperature (3) can be found from the formula

$$\frac{\Delta T_c}{T_c^{(0)}} \approx -\frac{2}{3}\frac{MT_c^{(0)}}{n}\left\langle\Delta\phi^2\right\rangle = -\frac{4\pi}{3}\frac{(MT_c^{(0)})^2}{n}4!\left\langle\frac{\Delta\phi^2}{u}\right\rangle a = -\frac{4\pi}{3}(2\pi)^2\frac{1}{[\zeta(3/2)]^{4/3}}4!\left\langle\frac{\Delta\phi^2}{u}\right\rangle an^{1/3}, \tag{7}$$

corresponding in Eq. (4) to

$$c_1 \approx -1103.09\left\langle\frac{\Delta\phi^2}{u}\right\rangle. \tag{8}$$

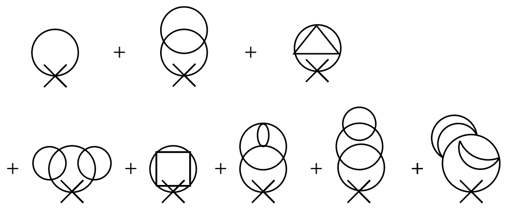

FIG. 1: Diagrams contributing to the expectation value $\langle\phi^2\rangle$.

The three-dimensional theory is superrenormalizable and requires only mass counterterms which shift the original bare mass $m$ to the renormalized mass $m_r$. A calculation of the Feynman diagrams in Fig. 1 yields the following five-loop perturbation expansion for the expectation value $\langle\phi^2/u\rangle$ [10, 11]

$$\begin{aligned}\left\langle\frac{\phi^2}{u}\right\rangle &= F(u) \equiv -\frac{N}{4\pi}\frac{m_r}{u} - a_2\frac{N\,(2+N)}{18\,(4\pi)^3}\frac{u}{m_r} + a_3\frac{N\,\left(16+10\,N+N^2\right)}{108\,(4\pi)^5}\left(\frac{u}{m_r}\right)^2 \\ &\quad - \left[\, a_{41}\frac{N(2+N)^2}{324\,(4\pi)^7} + a_{42}\frac{N\,\left(40+32\,N+8\,N^2+N^3\right)}{648\,(4\pi)^7} + a_{43}\frac{N\,\left(44+32\,N+5\,N^2\right)}{324\,(4\pi)^7}\right. \\ &\quad \left. + a_{44}\frac{N\,(2+N)^2}{324\,(4\pi)^7} + a_{45}\frac{N\,\left(44+32\,N+5\,N^2\right)\,u^4}{324\,m_r^3\,(4\pi)^7}\right]\left(\frac{u}{m_r}\right)^3 + \ldots.\end{aligned} \tag{9}$$

where $a_2 \equiv \log(4/3)/2 \approx 0.143841$ and the other constants are only known numerically [12]:

$$a_3 = 0.642144,\ a_{41} = -0.115069,\ a_{42} = 3.128107,\ a_{43} = 1.63,\ a_{44} = -0.624638,\ a_{45} = 2.39. \tag{10}$$

Writing the above expansion up to the $L$th term as $F_L(u) = \Sigma_{l=-1}^{L} f_l(u/4\pi m_r)^l$, the expansion coefficients for the relevant number of components $N = 2$ are [12]:

$$f_{-1} = -126.651\times10^{-4},\quad f_0 = 0,\quad f_1 = -4.04837\times10^{-4},\quad f_2 = 2.39701\times10^{-4},\quad f_3 = -1.80\times10^{-4}. \tag{11}$$

We need the value of the series $F_L(u)$ in the critical limit $m_r \to 0$, which is obviously equivalent to the strong-coupling limit of $F_L(u)$. As mentioned above, this limit should be most accurately found with the help of variational perturbation theory [4, 8, 9].

If the series were of quantum mechanical origin, we could find this limit by applying the rules of *naive* variational perturbation theory [13]. We form the sequence of truncated expansions $F_L(u)$ for $1, 2, 3$ and replace each term $(u/m_r)^l$ by $K^l[1-1]_{L-l}^{-l/2}$ where the symbol $[1-1]_k^r$ is defined as the binomial expansion of $(1-1)^r$ truncated after the $k$th term

$$[1-1]_k^r \equiv \sum_{i=0}^{k}\binom{r}{i}(-1)^i = (-1)^k\binom{r-1}{k}. \tag{12}$$

TABLE I: Trial functions for the naive quantum-mechanical variational perturbation expansion

$W_1^{\rm QM} = -0.0596831K^{-1} - 0.0000322159\,K,$

$W_2^{\rm QM} = -0.0497359K^{-1} - 0.0000483239\,K + 1.51792\,10^{-6}\,K^2,$

$W_3^{\rm QM} = -0.0435189K^{-1} - 0.0000604049\,K + 3.03584\,10^{-6}\,K^2 - .908\,10^{-7}\,K^3.$

The resulting expressions must be optimized in the variational parameter $K$. They are listed in Table I. The approximants $W_{1,2,3}^{\rm QM}$ have extrema $W_{1,2,3}^{\rm QMext} \approx -0.00277$, $+0.00405$, $-0.0029$, corresponding, via (8), to $c_1 \approx 3.059, -4.46, 3.01$. These values have previously been obtained in Ref. [10] in a much more complicated way via a so-called $\delta$-expansion. Note the negative sign of the second approximation arising from the fact that an extremum exists only at negative $K$. According to our rules of variational perturbation theory one should, in this case, use the saddle point at positive $K$ which would yield $W_2^{\rm QM} = -0.00153$ corresponding to $c_1 \approx 1.69$ rather than -4.46, leading to the more reaonable approximation sequence $c_1 \approx 3.059$, 1.69, 3.01, which shows no sign of convergence. In $W_3^{\rm QM}$, there is also a pair of complex extrema from which the authors of Ref. [10] extract the real part Re $\tilde W_{3\,{\rm complex}}^{\rm QM} \approx -0.00134$ corresponding to $c_1 \approx 1.48$, which they state as their final result. There is, however, no acceptable theoretical justification for such a choice [14].

This lack of convergence is not astonishing since naive quantum mechanical variational perturbation theory is inapplicable to field theory, contrary to ubiquitous statements in the literature [15]. A simple but essential modification is necessary to allow for the well-known fact that there are *anomalous dimensions* in the critical regime of fluctuating fields. This modification was discovered in Ref. [8] and tested by the fact it reproduces in $D = 4 - \epsilon$ dimensions *exactly* the known $\epsilon$ expansions of renormalization group theory [9]. In $D = 3$ dimensions, it leads to the most accurate critical exponents so far (see in particular Chapters 19 and 20 in the textbook [4]).

The correct procedure goes as follows: We form the logarithmic derivative of the expansion (9):

$$\beta\,(u) \equiv \frac{\partial \log F(u)}{\partial \log u} = -1 + 2\frac{f_1}{f_{-1}}\left(\frac{u}{m_r}\right)^2 + 3\frac{f_2}{f_{-1}}\left(\frac{u}{m_r}\right)^3 + \left(4\frac{f_3}{f_{-1}} - 2\frac{f_1^2}{f_{-1}^2}\right)\left(\frac{u}{m_r}\right)^4 + \dots\,. \tag{13}$$

In order for $F(u)$ to go to a constant in the critical limit $m_r \to 0$, this function must go to zero in the strong-coupling limit $u \to \infty$. Writing the expansion as $\beta_L\,(u) = -1 + \Sigma_{l=2}^L\, b_l (u/4\pi m_r)^l$, the coefficients are

$$b_2 = 0.0639293, \quad b_3 = -0.056778, \quad b_4 = 0.0548799. \tag{14}$$

The sums $\beta_L(u)$ have to be evaluated for $u \to \infty$ allowing for the universal anomalous dimension $\omega$ by which the physical observables of $\phi^4$-theories approach the scaling limit [3, 4]. This $\omega$ coincides with the famous Wegner exponent [5] of renormalization group theory. Here it appears in the variational expression for the strong-coupling limit which is found [8, 9] by replacing $(u/m_r)^l$ by $K^l[1-1]_{L-l}^{-ql/2}$, where $q \equiv 2/\omega$. Thus we obtain the variational expressions

$$W_3^\beta = -1 + \left(\frac{2\,f_1}{f_{-1}} + \frac{2\,f_1\,q}{f_{-1}}\right)K^2 + \frac{3\,f_2}{f_{-1}}K^3 \tag{15}$$

$$W_4^\beta = -1 + \left(\frac{2\,f_1}{f_{-1}} + \frac{3\,f_1\,q}{f_{-1}} + \frac{f_1\,q^2}{f_{-1}}\right)K^2 + \left(\frac{3\,f_2}{f_{-1}} + \frac{9\,f_2\,q}{2\,f_{-1}}\right)K^3 + \left(\frac{-2\,{f_1}^2}{{f_{-1}}^2} + \frac{4\,f_3}{f_{-1}}\right)K^4 \tag{16}$$

The first has a vanishing extremum at $\omega_3 = 0.592$, the second has neither an extremum nor a saddle point. However, a complex pair of extrema lies reasonably close to the real axis at $\omega_4 = 0.635 \pm 0.116$, whose real part is not far from the true exponent of approach $\omega_\infty \approx 0.80$ [3, 4], to which $\omega_L$ will converge for order $L \to \infty$ [8]. Given these $\omega$-values, we now form the variational expressions $W_L$ from $F_L$ by the replacement $(u/m_r)^l \to K^l[1-1]_{L-l}^{-ql/2}$, which are

$$W_2 = f_{-1}\left(1 - \frac{3}{4}q + \frac{1}{8}q^2\right)K^{-1} + f_1 K, \tag{17}$$

$$W_3 = f_{-1}\left(1 - \frac{11}{13}q + \frac{1}{4}q^2 - \frac{1}{48}q^3\right)K^{-1} + f_1\left(1 + \frac{q}{2}\right)K + f_2 K^2, \tag{18}$$

$$W_4 = f_{-1}\left(1 - \frac{25}{24}q + \frac{35}{96}q^2 - \frac{5}{96}q^3 + \frac{1}{384}q^4\right)K^{-1} + f_1\left(1 + \frac{3}{4}q + \frac{1}{8}q^2\right)K + f_2(1+q)K^2 + f_2 K^3. \tag{19}$$

The lowest function $W_2$ is optimized with the naive growth parameter $q = 1$ since to this order no anomalous value can be deternined from the zero of the beta function (13). The optimal result is $W_2^{\rm opt} = -\sqrt{\log[4/3]/6}/8\pi^2 \approx -0.00277$ corresponding to $c_1 \equiv 3.06$. The next function $W_3$ is optimized with the above determined $q_3 = 2/\omega_3$ and yields $W_3^{\rm opt} \approx -0.00096$ corresponding to $c_1 \equiv 1.06$. Although $\omega_4$ is not real we shall insert its real part into $W_4$ and find $W_4^{\rm opt} \equiv -0.00096$ corresponding to $c_1 \equiv 1.06$. The three values of $c_1$ for $\bar{L} \equiv L - 1 = 1, 2, 3$ can well be fitted by a function $c_1 \approx 1.053 + 2/\bar{L}^6$ (see Fig. 2). Such a fit is suggested by the general large-$L$ behavior $a + be^{-c\bar{L}^{1-\omega}}$ which was derived in Refs. [13]. Due to the smallness of $1 - \omega \approx 0.2$, this can be replaced by $\approx a' + b'/\bar{L}^s$.

Alternatively, we may optimize the functions $W_{1,2,3}$ using the known precise value of $q_\infty = 2/\omega_\infty \approx 2/0.8$. Then $W_2$ turns out to have no optimum, whereas the others yield $W_{3,4}^{\rm opt} \approx -0.000529, -0.000706$, corresponding via Eq. (8) to $c_1 = 0.584, 0.78$. If these two values are fitted by the same inverse power of $\bar{L}$, we find $c_1 \approx 0.798 - 13.7/\bar{L}^6$.

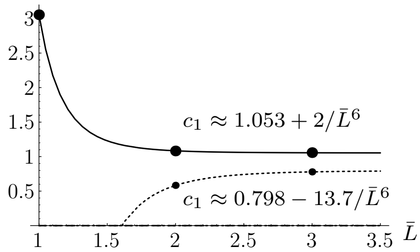


FIG. 2: The three approximants for $c_1$ plotted against the order of variational approximation $\bar{L} \equiv L - 1 = 1, 2, 3$, and extrapolation to the infinite-order limit.

From the extrapolations to infinite order we estimate $c_{1,\infty} \approx 0.93 \pm 0.13$, such that the critical temperature shift is

$$\frac{\Delta T_c}{T_c^{(0)}} \approx (0.93 \pm 0.13)\, a n^{1/3}. \tag{20}$$

Our result (20) is to be compared with latest Monte Carlo data which estimate $c_1 \approx 1.32 \pm 0.02$ [16, 17]. Other analytic estimates are $c_1 \approx 2.90$ [18], 2.33 from a $1/N$-expansion [19]), 1.71 from a next-to-leading order in a $1/N$-expansion [20], and 3.059 from an inapplicable $\delta$-expansion [21] to three loops, and 1.48 from the same $\delta$-expansion to five loops, with a questionable evaluation at a complex extremum [10], and some wrong expansion coefficients (see [12]). Remarkably, our result lies close to the average between the latest and the first Monte Carlo result $c_1 \approx 0.34 \pm 0.03$ in Ref. [22].

As a cross check of the reliability of our theory consider the result in the limit $N \to \infty$. Here we drop the first term in the expansion (9) which vanishes at the critical point (but would diverge for $N \to \infty$ at finite $m_r$). The remaining expansion coefficients of $\langle \phi^2/u \rangle /N$ in powers of $Nu/4\pi m_r$ are

$$f_1 = -6.35917\, 10^{-4}, \quad f_2 = 4.7315\, 10^{-4}, \quad f_3 = -3.84146\, 10^{-4}. \tag{21}$$

Using the $N \to \infty$ limit of $\omega$ which is equal to 1 implying $q = 2$ in Eqs. (18) and (19), we obtain the two variational approximations

$$W_2^\infty = -0.00127183K + 0.00047315K^2, \quad W_3^\infty = -0.00190775K + 0.00141945K^2 - 0.000384146K^3, \tag{22}$$

whose optima yield the approximations $c_1 \approx 1.886$ and 2.017, converging rapidly towards the exact large-$N$ result 2.33 of Ref. [19], with a 10% error as in the $N = 2$-result (20).

Numerically, the first two $1/N$-corrections found from a fit to large-$N$ results obtained by using the known large-$N$ expression for $\omega = 1 - 8(8/3\pi^2 N) + 2(104/3 - 9\pi^2/2)(8/3\pi^2 N)^2$ [23] produce a finite-$N$ correction factor $(1 - 3.1/N + 30.3/N^2 + \ldots)$, to be compared with $(1 - 0.527/N + \ldots)$ obtained in Ref. [20].

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integrals and confirming the correctness of the coefficients (10) of Ref. [11] rather than those of Ref. [10] listed in [12].

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