Physical mechanisms generating
spontaneous symmetry breaking
and a hierarchy of scales

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Abstract

We discuss the phase transition in $3 + 1$ dimensional $\lambda\Phi^4$ theory from a very physical perspective. The particles of the symmetric phase ('phions') interact via a hard-core repulsion and an induced, long-range $-1/r^3$ attraction. If the phion mass is sufficiently small, the lowest-energy state is not the ‘empty’ state with no phions, but is a state with a non-zero density of phions Bose-Einstein condensed in the zero-momentum mode. The condensate corresponds to the spontaneous-symmetry-breaking vacuum with $\langle \Phi \rangle \neq 0$ and its excitations (“phonons” in atomic-physics language) correspond to Higgs particles. The phase transition happens when the phion’s physical mass $m$ is still positive; it does not wait until $m^2$ passes through zero and becomes negative. However, at and near the phase transition, $m$ is much, much less than the Higgs mass $M_h$. This interesting physics coexists with “triviality;” all scattering amplitudes vanish in the continuum limit, but the vacuum condensate becomes infinitely dense. The ratio $m/M_h$, which goes to zero in the continuum limit, can be viewed as a measure of non-locality in the regularized theory. An intricate hierarchy of length scales naturally arises. We speculate about the possible implications of these ideas for gravity and inflation.
1. Introduction

Spontaneous symmetry breaking is an essential component of current theories of particle physics. All particles in the Standard Model acquire their masses from a non-vanishing expectation value \(\langle \Phi \rangle \neq 0\) of a self-interacting scalar field. The idea is simple and has a long history, so one might think that little remains to be understood. However, a basic question remains to be settled: the nature of the phase transition in the \(\lambda \Phi^4\) scalar field theory.

At the classical level one need only look at the potential

\[
V_{\text{cl}}(\phi) = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4
\]

(1.1)

to see that the phase transition, as one varies the \(m^2\) parameter, is second order and occurs at \(m^2 = 0\). In the quantum theory, however, the question is more subtle. Clearly, the symmetric vacuum is locally stable if its excitations have a physical mass \(m^2 > 0\) and locally unstable if \(m^2 < 0\). However, there remains the question of whether an \(m^2 > 0\) symmetric vacuum is necessarily globally stable. Could the phase transition actually be first order, occurring at some small but positive \(m^2\)?

The standard approximation methods for the quantum effective potential \(V_{\text{eff}}(\phi)\) give contradictory results on this crucial issue [1]. The straightforward one-loop approximation predicts a first-order transition occurring at a small but positive value of the physical (renormalized) mass squared, \(m^2 = m^2_c > 0\), so that the \(m^2 = 0\) case lies within the broken phase. On the other hand, the “renormalization-group-improved” result obtained by resumming the leading-logarithmic terms [1] predicts a second-order transition at \(m^2 = 0\). The conventional view is that the latter result is trustworthy while the former is not. The argument is that, for \(0 < m^2 < m^2_c\), the one-loop potential’s non-trivial minimum occurs only where the one-loop “correction” term is as large as the tree-level term. However, there is an equally strong reason to distrust the “RG-improved” result in the same region of \(m^2\) and \(\phi\); it amounts to re-summing a geometric series of leading logs that is actually a divergent series [2]. Moreover, the qualitative disagreement arises from a change in \(V_{\text{eff}}\) that in the crucial region is quantitatively tiny — exponentially small in the coupling constant. One cannot trust perturbation theory, improved or otherwise, at that level. Thus, in \(\lambda \Phi^4\) theory [3] it is unsafe to draw any firm conclusion from either method; other approaches must be sought.

The Gaussian approximation [4] provides a clue. In 3+1 dimensions it produces a result in agreement with the one-loop effective potential [5]. This is not because it contains no
non-vanishing corrections beyond the one-loop level; it does, but those terms do not alter the functional form of the result. When reparametrized in terms of a physical mass and field, the renormalized result is exactly the same [6].

The continuum limit of \( \lambda \Phi^4 \) in 4 space-time dimensions is almost certainly ‘trivial’ [7, 8, 9]. Thus, a key consideration is what ‘triviality’ implies about the effective potential. Initially, one might presume that ‘triviality’ implies a quadratic effective potential, as in free-field theory. However, that presumption accords with none of the approximate methods and is far from being satisfactory [10]. Instead, we advocate the following viewpoint [11, 2, 12]: if a theory is ‘trivial,’ then its effective potential should be physically indistinguishable from the classical potential plus a zero-point-energy contribution of free-field form arising from fluctuations:

\[
V_{\text{triv}}(\phi) \equiv V_{\text{cl}}(\phi) + \frac{1}{\sqrt{\lambda}} \sum_k \frac{1}{2} \sqrt{k^2 + M_k^2(\phi)}.
\]  

(1.2)

Here \( M(\phi) \) denotes the mass of the shifted (‘Higgs’) field \( h(x) = \Phi(x) - \phi \) in the presence of a background field \( \phi \). After mass renormalization and subtraction of a constant term, \( V_{\text{triv}}(\phi) \) consists of \( \phi^2, \phi^4, \) and \( \phi^4 \ln \phi^2 \) terms. Any detectable difference from this form would imply interactions of the \( h(x) \) field — and there are none if the theory is ‘trivial.’ In other words, ‘triviality’ implies that the exact result for the continuum-limit effective potential should be physically indistinguishable from the one-loop result. Notice that we say “physically indistinguishable from …” not “equal to …;” it is not that multi-loop graphs produce no contributions but that those contributions affect both the effective potential and \( M_k^2(\phi) \) in a way that preserves the functional form implied by (1.2), up to terms that vanish in the continuum limit.

This viewpoint explains the exact agreement between the Gaussian and one-loop results noted above. Moreover, it implies that there is an infinite class of “triviality-compatible” approximations, all yielding the same result. Such approximations can be arbitrarily complex provided they have a variational or CJT structure [13], with the shifted ‘Higgs’ field \( h(x) = \Phi(x) - \phi \) having a propagator determined variationally by solving a non-perturbative gap-equation. If the approximation is “triviality compatible” then this propagator reduces to a free-field propagator in the infinite-cutoff limit. In that limit all differences among these various approximations can be absorbed into a redefinition of \( \lambda \), which makes no difference when the effective potential is expressed in terms of physical renormalized quantities [8, 14]. (An explicit example of such a calculation, beyond the Gaussian approximation, is provided by ref [13].)
The form $V_{\text{triv}}(\phi)$ differs from the prediction of renormalization-group-improved perturbation theory (RGIPT). Many readers, we realize, will balk at any criticism of RGIPT [16]. However, this issue can be addressed objectively; for instance, the two predictions can be tested against a sufficiently precise lattice Monte-Carlo calculation [14, 17]. Data from such a lattice simulation [17] support our position; an excellent fit ($\chi^2/\text{d.o.f.} \leq 1$) is obtained with $V_{\text{triv}}$, whereas the form predicted by RGIPT is unable to fit the data ($\chi^2/\text{d.o.f.} \sim 6 - 10$). This evidence certainly justifies us in pursuing our picture further.

In some respects the differences between the RGIPT and $V_{\text{triv}}$ forms are quite small and subtle. The following toy model helps to illustrate this point: Consider a potential

$$V_{\text{toy}} = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4 \left(1 + \epsilon \ln \frac{\phi^2}{\mu^2}\right),$$

(1.3)

where $\mu$ is some mass scale and $\epsilon$ is a small parameter. (The real case is like $\epsilon \propto \lambda$, modulo some technicalities, but in this toy model we treat $\epsilon$ as a separate parameter.) For $\epsilon = 0$ one has a second-order phase transition, occurring at $m^2 = 0$, as one varies the $m^2$ parameter. However, for any positive $\epsilon$, no matter how small, one has a first-order transition, occurring at a positive $m^2$. The size of $m^2$ involved is exponentially small in units of $\mu^2$, $O(\lambda \epsilon e^{-1/\epsilon}) \mu^2$. The vacuum value of $\phi^2$ is also exponentially small, $O(\epsilon^{-1/\epsilon}) \mu^2$. The difference between $V_{\text{toy}}$ and $V_{\text{cl}}$ in this region of $\phi$ is only $O(\lambda \epsilon e^{-2/\epsilon}) \mu^4$. This toy model illustrates the point that a very weak first-order phase transition becomes indistinguishable from a second-order transition if one does not look on a fine enough scale. If one varies $m^2$ on a scale of $\mu^2$ one sees what looks like a second-order transition. Only if one varies $m^2$ on a much finer scale does one see that the transition is first order, exhibiting a small but non-zero jump in the order parameter, and occurs at a small but non-zero $m^2$.

We can now formulate the specific puzzle addressed in this paper. Suppose that spontaneous symmetry breaking does indeed coexist with a physical mass $m^2 \geq 0$ for the excitations of the symmetric phase. Those excitations would then be real particles — as real as electrons or quarks (though, like quarks, they would not be directly observable); for brevity we call them ‘phions.’ The puzzle is this: How is it possible for the broken-symmetry vacuum — a condensate with a non-zero density of phions — to have a lower energy density than the ‘empty’ state with no phions? The $\lambda \Phi^4$ interaction corresponds to a repulsive ‘contact’ interaction between phions [18] and one would think that any state made out of positive-mass particles with a repulsive interaction would necessarily have a positive energy density.

The solution to this puzzle is the realization that the phion-phion interaction is not
always repulsive; there is an induced interaction that is attractive. Moreover, as \( m \to 0 \) the attraction becomes so long range, \(-1/r^3\), that it generates an infrared-divergent scattering length. As we shall see, this long-range attraction makes it energetically favourable for the condensate to form spontaneously. This leads to a simple picture — a physical mechanism — for spontaneous symmetry breaking.

The physics is directly related to the Bose-Einstein (BE) condensation of a dilute, non-ideal, Bose gas (a phenomenon observed recently in atom-trap experiments [19]). The theory for this is very well established [20, 21, 22]. The elementary excitations of an atomic condensate represent not single-atom motions but collective motions — quantized pressure waves, or “phonons.” In this language the Higgs particle is the phonon excitation of the phion condensate.

One might then ask: . . . but how is this interesting physics consistent with ‘triviality’? The interaction between phions should vanish in the infinite-cutoff limit (corresponding to shrinking the intrinsic phion size to zero). How can such a ‘trivial’ theory have a not-entirely-trivial ground state? The answer is that even an infinitesimal two-body interaction can induce a macroscopic change of the ground state if the vacuum contains an infinite density of condensed phions. Indeed we shall find that the condensate density is infinite in physical length units set by \( M_h^{-1} \), the inverse Higgs mass. Nevertheless, the condensate is infinitely dilute — the density is vanishingly small on a length scale set by the scattering length. This sort of subtlety reflects the existence of a hierarchy of scales. One length scale, the scattering length, vanishes (‘triviality’), while another, set by \( M_h^{-1} \), remains finite. In fact, an intricate hierarchy of scales emerges, as we discuss in Sect. 8.

In what follows we use units with \( \hbar = c = 1 \). For simplicity we consider the single-component \( \lambda \Phi^4 \) theory with a discrete reflection symmetry, \( \Phi \to -\Phi \). Since the field is Hermitian, the phion particle will be its own antiparticle. In Sect. 2 we discuss the interparticle potential between phions. Then in Sect. 3 we estimate the energy density of a phion condensate, with a given particle density, in a very simple and intuitive way. The result is confirmed in Sect. 4 by a calculation based on the Lee-Huang-Yang (LHY) treatment of a non-relativistic Bose gas. The resulting energy-density expression, equivalent to the field-theoretic effective potential, yields a phase transition which we analyze in Sect. 5. The excitations of the condensate are ‘phonons’ (Higgs particles) and in principle the physics can be described either in terms of phions or in terms of phonons. Sect. 6 discusses how the “renormalized field” associated with phonons is related to the original (“bare”) field associated with phions. In Sect. 7 the effective potential is written in manifestly finite form in terms of the renormalized field. Sect. 8 provides a brief summary.
and discusses the intricate hierarchy of length scales that arise. We conclude with some speculations about the possible wider implications of these ideas.

2. Interparticle potential between phions

In QED there is a well-known equivalence between the photon-exchange Feynman diagram and the Coulomb potential $1/r$. Similarly, pion exchange gives rise to the Yukawa potential $e^{-mr}/r$ in nuclear physics. The exchange of two massless neutrinos gives rise to a long-range $1/r^5$ potential [23]. In $\lambda\Phi^4$ theory it is well known that the fundamental interaction vertex corresponds to a $\delta^{(3)}(r)$ interaction [18]. However, as we now explain, the exchange of two virtual phions gives rise to an attractive long-range $-1/r^3$ interaction.

Consider the elastic collision of two particles of mass $m$ in the centre-of-mass frame. Let $q$ denote the 3-momentum transfer; let $\theta$ be the scattering angle; and let $E = \sqrt{p^2 + m^2}$ be the energy of each particle. A scattering matrix element $M$, obtained from Feynman diagrams, can be associated with an ‘equivalent interparticle potential’ that is basically the 3-dimensional Fourier transform of $M$:

$$V(r) = \frac{1}{4E^2} \frac{d^3q}{(2\pi)^3} e^{iq\cdot r} M.$$  \hspace{1cm} (2.1)

(For a detailed discussion see the review article of Feinberg et al [24].) This ‘equivalent potential’ is a function of the relative position $r$ (conjugate to $q$). In general it also depends parametrically on the energy, $E$, though this complication disappears in the non-relativistic limit, where $E \sim m$.

For the $\lambda\Phi^4$ theory the lowest-order Feynman diagram (see Fig. 1) gives $M_o = \lambda$ and the resulting potential is

$$V_o(r) = \frac{1}{4E^2} \lambda \delta^{(3)}(r).$$  \hspace{1cm} (2.2)

In a non-relativistic treatment of the theory this is the only interaction. The ‘triviality’ property is then reflected in the well-known fact that in quantum mechanics a 3-dimensional $\delta$-function interaction gives zero scattering amplitude [18].

Relativistically there are additional contributions to the equivalent interparticle potential, notably those produced by the one-loop ‘fish’ diagrams (see Fig. 2) corresponding to the three Mandelstam variables $s, t, u$ ($s = 4E^2$, $t = -q^2$, $s + t + u = 4m^2$). To evaluate these contributions, we first note that if $M$ depends only on $q \equiv |q|$, then Eq. (2.1) reduces to

$$V(r) = \frac{1}{4E^2} \frac{1}{(2\pi)^2} \int_0^\infty q^2 dq \frac{2\sin qr}{qr} M(q)$$  \hspace{1cm} (2.3)
This already shows that \( V(r) \) is spherically symmetric and naturally has a factor \( 1/r^3 \).

To evaluate the contribution from \( t \)-channel scattering, we begin with the case \( m = 0 \), where the matrix element is simple:

\[
\mathcal{M}_{t-\text{exch}}(q) = \frac{\lambda^2}{16\pi^2} \ln(q/\Lambda),
\]

where \( \Lambda \) is the ultraviolet cutoff. Substituting into Eq. (2.4) we find an integral that is not properly convergent but which can be made convergent by including a factor \( e^{-\epsilon y} \) and then taking the limit \( \epsilon \to 0 \) (physically, this corresponds to smearing out the point vertices). In this sense we have

\[
\int_0^\infty dy \, y \sin y = 0,
\]

and

\[
\int_0^\infty dy \, y \ln y \sin y = -\frac{\pi}{2}.
\]

The first equation implies that those terms independent of \( q \) in the matrix element do not give contributions to the potential for values of \( r \neq 0 \). Such terms, however, bring a contribution of the type \( \delta^{(3)}(r) \), as we see by returning to the form (2.1). The \( s \)-channel amplitude, for example, gives only a contribution of this type. These delta-function contributions can be absorbed into a redefinition of \( \lambda \), the strength of the repulsive potential in Eq. (2.2). In this way, we can include all possible diagrammatic contributions to the short-range repulsive interaction. Then, \( \lambda \) would become an effective parameter representing the actual physical strength of the repulsive contact interaction, rather than the bare coupling entering in the Lagrangian density.

Substituting Eq. (2.5) into (2.4) we find an attractive, long-range potential:

\[
V_{t-\text{exch}}(r) = -\frac{\lambda^2}{256\pi^4 E^2} \frac{1}{r^3}.
\]

An equal contribution is obtained from the \( u \)-channel diagram; in QM terms it corresponds to the amplitude \( f(\pi - \theta) \) that must be added to \( f(\theta) \) when dealing with identical-particle scattering. Note that, as a consequence of Eq. (2.6), there is no dependence on \( \Lambda \) in the \( -1/r^3 \) potential.

Taking into account the mass \( m \) of the exchanged particles yields the result (2.8) multiplied by a factor \( 2mrK_1(2mr) \), where \( K_1 \) is the modified Bessel function of order unity. This factor tends to unity as \( mr \to 0 \) and for large values of \( mr \) tends to \( \sqrt{\pi mr} e^{-2mr} \).
The exponential factor is like that of the Yukawa potential except that, since there are two exchanged particles, it is \( e^{-2mr} \) instead of \( e^{-mr} \). Physically, the \(-1/r^3\) potential arises from two short-range repulsive interactions linked by the quasi-free propagations of two virtual particles. (Exchanges of more than 2 particles over macroscopic distances would lead to contributions with a faster power-law fall-off.) Thus, higher-order contributions are accounted for by the same redefinition of \( \lambda \) mentioned above. If the short-range repulsive interaction has an actual strength \( \lambda \), then the \(-1/r^3\) attractive interaction is proportional to \( \lambda^2 \).

In summary: the interparticle potential is essentially given by the sum of a repulsive core, \( \delta^{(3)}(r) \), and an attractive term \(-1/r^3\) that is eventually cut off exponentially at distances greater than \(1/(2m)\). The long-range attraction between the phions has important effects, as we shall see in the next section.

3. Condensate energy density: a simple estimate

Consider a large number \( N \) of phions contained in a large box of volume \( V \). As in statistical mechanics, the thermodynamic limit requires \( N \to \infty \) and \( V \to \infty \) with the density \( n \equiv N/V \) being fixed. The 'empty' state corresponds to the special case \( n = 0 \). Since phions can be created and destroyed, the equilibrium value of \( n \) is to be determined by minimizing the energy density in the box. In this section we estimate the ground-state energy density for a given \( n \) in a very simple and intuitive way. Some tedious subtleties affecting numerical factors are ignored here. A proper calculation will be provided in the next section.

Assuming the density \( n \) is low, the relevant contributions to the total energy of the system are just the rest-masses \( Nm \) and the two-body interaction energies. Effects from three-body or multi-body interactions will be negligible provided the gas of phions is dilute. The two-body contribution is the number of pairs \((\frac{1}{2}N(N-1) \approx \frac{1}{2}N^2)\) multiplied by the average potential energy between a pair of phions:

\[
\bar{u} \sim \frac{1}{V} \int d^3r V(r). \tag{3.1}
\]

This averaging assumes that the particles are uniformly distributed over the box, which is valid since at zero temperature almost all the particles are condensed in the \( k = 0 \) mode. Thus, the total energy of the ground state is

\[
E_{\text{tot}} = Nm + \frac{1}{2}N^2\bar{u}, \tag{3.2}
\]
yielding an energy density
\[ \mathcal{E} \equiv \frac{E_{\text{tot}}}{V} = n m + \frac{\lambda}{2} n^2 \int d^3 r V(r). \] (3.3)

The potential \( V(r) \) consists of the \( \delta^{(3)}(r) \) term (2.2) and the \(-1/r^3\) term (2.8) (times 2 to include the \( u \)-channel). We may set \( E = m \) since almost all phions have \( k = 0 \). Thus, we find
\[ \mathcal{E} = n m + \frac{\lambda n^2}{8 m^2} - \frac{\lambda^2 n^2}{64 \pi^2 m^2} \int \frac{dr}{r}. \] (3.4)

The integral over \( r \) can be cut off at small \( r \) by introducing a ‘hard-core radius’ \( r_o \), corresponding to an ultraviolet regularization that smears out the \( \delta^{(3)}(r) \) point-like interaction. In addition, to avoid an infrared divergence, the integral must also have some large-distance cutoff, \( r_{\text{max}} \). Since, as noted in the last section, the phion mass introduces an exponential factor \( e^{-2m r} \) into the potential, we have \( r_{\text{max}} \leq 1/(2m) \). However, if \( m \) is very small another consideration is actually more important; namely, that the long-distance attraction between two phions becomes “screened” by other phions that interpose themselves. This immediately implies an \( r_{\text{max}} \) that depends on the density \( n \). In fact, \( r_{\text{max}} \) is naturally given by \( 1/(2M) \), where \( M \) is the mass of the quasiparticle excitations of the condensate, and it is easily seen that \( M^2 \) is proportional to \( n \) when \( m \) is small [24].

Hence, \( \mathcal{E} \) is given by a sum of \( n, n^2 \) and \( n^2 \ln n \) terms which represent, respectively, the rest-mass energy cost, the repulsion energy cost, and the energy gain from the long-range attraction. If the rest-mass \( m \) is small enough, then the \( n^2 \ln n \) term’s negative contribution can result in an energy density whose global minimum is not at \( n = 0 \) but at some specific, non-zero density \( n_v \). That is, even though the ‘empty’ state is locally stable, it can decay by spontaneously generating particles so as to fill the box with a dilute condensate of density \( n_v \).

The result can be translated into field-theory terms since the particle density \( n \) is proportional to the intensity of the field, \( \phi^2 \). In fact, as shown in the next section, one has \( n = \frac{\lambda}{2} m \phi^2 \). The energy density as a function of \( n \) then becomes the field-theoretic effective potential: \( \mathcal{E}(n) \equiv V_{\text{eff}}(\phi) \). (Of course this “potential” for the field is not to be confused with the “potential” between particles, \( V(r) \).) The estimate above leads to
\[ V_{\text{eff}}(\phi) = \frac{1}{2} m^2 \phi^2 + \frac{\lambda \phi^4}{32} - \frac{\lambda^2 \phi^4}{256 \pi^2} \ln \frac{r_{\text{max}}(\phi)}{r_o}. \] (3.5)

We can identify \( r_o \) with the reciprocal of an ultraviolet cutoff \( \Lambda \) and \( r_{\text{max}}(\phi) \) with \( 1/(2M) \), where \( M^2 \propto n \propto \phi^2 \). Thus, the essential form of the result is here. (The incorrect numerical factors could be straightened out with enough care, we believe.)
This simple approach gives some important insight. The reason there are no \( n^3, n^4 \ldots \) \((\phi^6, \phi^8, \ldots)\) terms is the diluteness of the gas. Furthermore, if the attractive potential had fallen off faster than \( 1/r^3 \) then \( \mathcal{E} \) would have had only \( n \) and \( n^2 \) terms. The crucial \( n^2 \ln n \) term arises from the infrared divergence of the integral in (3.4) which is tamed only by the screening effect of the background density \( n \). Thus the \( \phi^4 \ln \phi^2 \) term has two complementary interpretations: In field language it arises from the zero-point energy of the field fluctuations, while in particle language it arises from the long-range attraction between phions.

4. Condensate energy density: calculation à la LHY

In this section we compute the energy density using a relativistic version of the original Lee-Huang-Yang (LHY) analysis of Bose-Einstein condensation of a non-ideal gas \([20, 21]\). We emphasize that their analysis invokes neither a weak-coupling nor a semiclassical approximation. They appeal to two approximations: (1) low energy, so that the scattering is pure \( s \)-wave and (2) low density (diluteness). In terms of the phion-phion scattering length

\[
a = \frac{\lambda}{8\pi E}
\]  

(4.1) 

these conditions are:

'low-energy': \( ka \ll 1 \),  

(4.2)

'diluteness': \( na^3 \ll 1 \).  

(4.3)

Note that ‘low energy’ in the above sense does not imply ‘non-relativistic’ — because here (quite unlike the situation in atomic physics) we may have \( m \ll 1/a \).

We start from the \( \lambda \Phi^4 \) Hamiltonian:

\[
H = \int d^3x \left[ \frac{1}{2} \left( \Pi^2 + (\nabla \Phi)^2 + m^2 \Phi^2 \right) + \frac{\lambda}{4!} \Phi^4 \right] .  
\]  

(4.4)

The system is assumed to be contained within a finite box of volume \( V \) with periodic boundary conditions. There is then a discrete set of allowed modes \( \mathbf{k} \). In the end we will take the infinite-volume limit and the summation over allowed modes will go over to an integration: \( \sum_{\mathbf{k}} \rightarrow V \int d^3k/(2\pi)^3 \).

Annihilation and creation operators, \( a_{\mathbf{k}}, a_{\mathbf{k}}^\dagger \), are introduced through the plane-wave expansion

\[
\Phi(x, t) = \sum_{\mathbf{k}} \frac{1}{\sqrt{2V\omega_k}} \left[ a_{\mathbf{k}} e^{i\mathbf{k} \cdot \mathbf{x}} + a_{\mathbf{k}}^\dagger e^{-i\mathbf{k} \cdot \mathbf{x}} \right] ,  
\]  

(4.5)
where \( E_k = \sqrt{k^2 + m^2} \). The \( a_k \)'s are time dependent (in the free-field case they would be proportional to \( e^{-iE_k t} \)) and satisfy the commutation relations

\[
[a_k, a_{k'}^\dagger] = \delta_{k,k'}.
\] (4.6)

The Hamiltonian includes “normal ordering” symbols \( \ldots \) so that the quadratic part of \( H \) is just

\[
H_2 = \sum_k E_k a_k^\dagger a_k.
\] (4.7)

For comparison with the non-relativistic calculation, it is convenient to subtract from the Hamiltonian a term \( \mu c \hat{N} \), where \( \hat{N} \) is the operator that counts the number of phions:

\[
\hat{N} = \sum_k a_k^\dagger a_k.
\] (4.8)

At the end we shall set the chemical potential \( \mu = 0 \). However, in the non-relativistic context one should put \( \mu = m \) to take into account that, then, the rest-mass energy is not counted as part of a particle’s kinetic energy. Therefore, the correct definition of the total energy of the system, in the non-relativistic context, is obtained by subtracting \( m \) for each particle, so that

\[
H_{NR} = H - m \hat{N}.
\]

When a system of \( N \) bosons undergoes Bose-Einstein condensation, then the lowest energy mode becomes macroscopically populated, below some critical temperature. That is, there are \( N_0 \) particles in the \( k = 0 \) mode, with \( N_0 \) being a finite fraction of the total number \( N \). At zero temperature, if the gas is dilute, almost all the particles are in the condensate; \( N_0(T = 0) \sim N \). In fact, the fraction not in the condensate is of order \( \sqrt{na_3} \) \( ^{20,21} \) and so is negligible in the dilute approximation. We then have \( a_0^\dagger a_0 \sim N \), and so we can consider \( a_0 \) to be essentially the c-number, \( \sqrt{N} \). (Of course \( a_0 \) still has an operator part, but any relevant matrix elements of this part are only of order unity, negligible in comparison to the c-number part \( \sqrt{N} \).) From the expansion (4.4) we then get the expectation value

\[
\phi = \langle \Phi \rangle = \frac{1}{\sqrt{2Vm}}(a_0 + a_0^\dagger) = \sqrt{\frac{2N}{Vm}}.
\] (4.9)

Hence, the particle density \( n \equiv N/V \) is given by

\[
n = \frac{\phi^2}{2m}.
\] (4.10)

as anticipated in the previous section. With this identification, setting \( a_0 = a_0^\dagger = \sqrt{N} \) is equivalent to shifting the quantum field \( \Phi \) by a constant term \( \phi \).
Making this substitution yields
\[
H_{\text{eff}} - \mu_c \hat{N} = V \left[ (m - \mu_c) n + \frac{\lambda n^2}{6m^2} \right] + \sum_{k \neq 0} \left[ a_k^\dagger a_k \left( E_k - \mu_c + \frac{\chi}{2E_k} \right) + \frac{\chi}{4E_k} \left( a_k a_{-k} + a_k^\dagger a_{-k}^\dagger \right) \right]
\]
(4.11)
with
\[
\chi = \frac{\lambda n}{m} = \frac{1}{2} \lambda \phi^2.
\]
(4.12)

As stressed in Ref. [28], this result contains all interactions of condensate particles between themselves and all interactions between condensate and non-condensate particles. It neglects interactions among the non-condensate particles, which is justified because there are so few of them; their density is smaller than \( n \) by a factor of \( \sqrt{n a^2} \) [21]. We stress that the justification here is not weakness of interaction but scarcity of interactors; i.e., low density and not weak coupling.

To diagonalize the Hamiltonian (4.11) we can define new annihilation and creation operators \( b_k, b_k^\dagger \) (for \( k \neq 0 \)). The linear transformation
\[
a_k = \frac{1}{\sqrt{1 - \alpha_k^2}} \left( b_k - \alpha_k b_{-k}^\dagger \right)
\]
(4.13)
and its Hermitian conjugate are called the Bogoliubov transformation. The quanta annihilated and created by the operators \( b_k, b_k^\dagger \) are called ‘phonons’ or ‘quasiparticles’ to distinguish them from the ‘particles’ associated with the original operators \( a_k, a_k^\dagger \). The function \( \alpha_k \) is fixed by the requirement that in the Hamiltonian Eq. (4.11), the coefficients of the \( b_k b_{-k} \) and \( \tilde{b}_k b_{-k}^\dagger \) terms vanish. This fixes
\[
\alpha_k = 1 + x^2 - x \sqrt{x^2 + 2}, \quad x^2 \equiv \frac{2}{\chi} E_k (E_k - \mu_c).
\]
(4.14)
The result then takes the form
\[
H_{\text{eff}} - \mu_c \hat{N} = E_{\text{tot}} + \sum_{k \neq 0} \tilde{E}_k b_k^\dagger b_k.
\]
(4.15)
Apart from the constant term \( E_{\text{tot}} \), which we discuss below, this is analogous to Eq. (4.7) but with a different spectrum (energy-momentum relationship) for the quasiparticles:
\[
\tilde{E}_k = \left( \frac{1 + \alpha_k}{1 - \alpha_k} \right) E_k,
\]
(4.16)
which yields
\[
\tilde{E}_k = (E_k - \mu_c) \sqrt{1 + \frac{\chi}{E_k (E_k - \mu_c)}}.
\]
(4.17)
In the non-relativistic limit, setting $\mu_c = m$ and $E_k = \sqrt{k^2 + m^2} \approx m + k^2/(2m) + \ldots$, one obtains the famous ‘Bogoliubov spectrum’:

$$\tilde{E}_k^{\text{NR}} = \frac{k}{2m} \sqrt{k^2 + 2\chi},$$

(4.18)

with its characteristic linear behaviour, as $k \to 0$, for the ‘phonon’ excitations of a dilute Bose gas at low temperature. In the relativistic case, where $\mu_c = 0$, one has instead

$$\tilde{E}_k = \sqrt{E_k^2 + \chi} = \sqrt{k^2 + m^2 + \chi}.$$  

(4.19)

This has the normal form for a relativistic energy-momentum relation and we can identify the mass of the ‘quasiparticle’ excitations as:

$$M^2(\phi) = m^2 + \chi = m^2 + \frac{1}{2} \lambda \phi^2.$$  

(4.20)

The constant term in Eq. (4.15) above is given (for $\mu_c = 0$) by

$$E_{\text{tot}} = \sqrt{nm + \frac{\lambda n^2}{6m^2}} - \frac{\chi}{4} \sum_{k \neq 0} \frac{\alpha_k}{E_k},$$

(4.21)

The term in square brackets is just the ‘classical’ energy density

$$V_{\text{cl}}(\phi) = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4$$

(4.22)

as sees by substituting $n = \frac{1}{2} m \phi^2$. The last term in (4.21) arises because $b_k b_k^\dagger = b_k^\dagger b_k + 1$. To evaluate it we substitute for $\alpha_k$ from Eq. (4.14), and use $x \sqrt{x^2 + 2} = 4E_k \tilde{E}_k / (\lambda \phi^2)$ to obtain

$$V_{\text{eff}} = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4 + I_1(M) - I_1(m) - \frac{1}{4} \lambda \phi^2 I_0(m),$$

(4.23)

where

$$I_1(M) \equiv \int \frac{d^3k}{(2\pi)^3} \frac{1}{2} \tilde{E}_k, \quad I_1(m) \equiv \int \frac{d^3k}{(2\pi)^3} \frac{1}{2} \tilde{E}_k, \quad I_0(m) \equiv \int \frac{d^3k}{(2\pi)^3} \frac{1}{2} \tilde{E}_k.$$  

(4.24)

In field-theory language $I_1(M)$ represents the zero-point fluctuations of a free scalar field of mass $M = M(\phi)$, and the last two terms of Eq.(4.23) represent the subtractions associated with the normal ordering of the Hamiltonian (4.4). Such subtractions remove the quartic divergence $\sim \Lambda^4$ and the quadratic divergence $\sim \Lambda^2$ that are contained in $I_1(M)$, leaving only a logarithmic divergence $\ln \Lambda$. Explicit calculation gives

$$V_{\text{eff}}(\phi) = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4 + \frac{\lambda^2}{256\pi^2} \phi^4 \left[ \ln(\frac{1}{4} \lambda \phi^2 / \Lambda^2) - \frac{1}{2} + F\left(\frac{m^2}{\frac{1}{4} \lambda \phi^2}\right) \right],$$

(4.25)
where
\[ F(y) = \ln(1 + y) + \frac{y(4 + 3y)}{2(1 + y)^2}, \quad y \equiv \frac{m^2}{\lambda \phi^2}. \] (4.26)

The result (4.25) coincides with the famous one-loop result [1]. However, our point is that the result is justified by the ‘low-energy’ and ‘diluteness’ assumptions, without appealing to perturbative or semiclassical approximations. This point is well known in the non-relativistic case [21, 28]. For small \( m \), when the \( F(y) \) term can be neglected, the result (4.25) has the same structure found in the intuitive calculation in the preceding section, Eq. (4.3).

5. The phase transition

We now analyze the energy-density expression Eq. (4.25) to find where the phase transition occurs. It is easy to guess that \( m^2 \) will have to be small, so we expect \( y \equiv m^2/(\lambda \phi^2) \ll 1 \) everywhere except very near the origin. This implies that the phonon mass \( M^2(\phi) \), Eq. (4.20), will be much larger than \( m^2 \) and will become essentially \( \frac{1}{2} \lambda \phi^2 \). In this regime the mass will be relevant only in the rest-mass energy term \( \frac{1}{2} m^2 \phi^2 \) and we may neglect the \( F(y) \) term. We proceed to do so, but we shall return at the end to verify that this is justified. Thus, we start from

\[ V_{\text{eff}} = \frac{1}{2} m^2 \phi^2 + \frac{\lambda}{4!} \phi^4 + \frac{\lambda^2}{256 \pi^2} \phi^4 \left[ \ln\left( \frac{\lambda \phi^2}{\Lambda^2} \right) - \frac{1}{2} \right], \] (5.1)

where we have added a ‘\( B \)’ subscript to \( \phi \) to emphasize that it is a ‘bare’ (unrenormalized) field. If \( V_{\text{eff}} \) has a pair of extrema \( \phi_B = \pm v_B \), then \( v_B \) is a solution of \( dV_{\text{eff}}/d\phi_B = 0 \), which gives

\[ m^2 + \frac{\lambda}{6} v_B^2 + \frac{\lambda^2}{64 \pi^2} v_B^2 \ln\left( \frac{\lambda v_B^2}{\Lambda^2} \right) = 0. \] (5.2)

This condition allows us to eliminate \( \Lambda \) in favour of \( v_B \), so that the effective potential (5.1) can be expressed equivalently as

\[ V_{\text{eff}} = \frac{1}{2} m^2 \phi_B^2 (1 - \frac{\phi_B^2}{2 v_B^2}) + \frac{\lambda^2}{256 \pi^2} \phi_B^4 \left[ \ln\left( \frac{\phi_B^2}{v_B^2} \right) - \frac{1}{2} \right]. \] (5.3)

We denote by \( v_0 \) the value of \( v_B \) in the case \( m^2 = 0 \); it is given by

\[ v_0^2 = \frac{2 \Lambda^2}{\lambda} \exp\left( \frac{-32 \pi^2}{3 \lambda} \right). \] (5.4)

The original equation (5.2) can then be re-written as:

\[ f(v_B^2) \equiv -\frac{\lambda^2}{64 \pi^2} v_B^2 \ln(v_B^2/v_0^2) = m^2. \] (5.5)
A graph of \( f(v_B^2) \) starts from zero at \( v_B^2 = 0 \), reaches a maximum at \( v_B^2 = e^{-1}v_0^2 \) and then decreases, becoming negative when \( v_B^2 > v_0^2 \). Equating this to \( m^2 \) we see that: (i) if \( m^2 \) is positive and larger than the maximum value of \( f \) then Eq. (5.3) has no real roots. In this case \( V_{\text{eff}} \) has a single minimum located at \( \phi_B = 0 \). (ii) if \( m^2 \) is positive but not too large, then Eq. (5.3) has two roots. In that case \( V_{\text{eff}} \) has a local minimum at \( \phi_B = 0 \), then a maximum (at the smaller \( v_B \) root) and then a minimum (at the larger root, the true \( v_B \)). (iii) if \( m^2 \) is negative there is a unique root, with \( v_B^2 \) greater than \( v_0^2 \). In that case, the origin \( \phi_B = 0 \) is a maximum of \( V_{\text{eff}} \) and \( v_B \) is an absolute minimum.

Case (i) is not very interesting since it does not show condensation and spontaneous symmetry breaking. Case (iii) shows spontaneous symmetry breaking but, with \( m^2 \) negative, the phions would not be particles in the ordinary sense. Our interest in this paper is with case (ii) where \( m^2 \) is positive but less than the maximum of the function \( f \):

\[
m^2 < \frac{\lambda^2}{64\pi^2} \frac{v_0^2}{e}. \tag{5.6}
\]

This condition ensures that non-trivial minima of \( V_{\text{eff}} \) exist. A stronger bound on \( m^2 \) is needed if the minimum at \( v_B \) is to have lower energy density than the symmetric vacuum. For this we need

\[
V_{\text{eff}}(\phi_B = \pm v_B) = \frac{1}{4}(m^2 - \frac{\lambda^2}{128\pi^2}v_B^2)v_B^2 \tag{5.7}
\]

to be negative. Combined with Eq. (5.3), this gives

\[
m^2 \leq \frac{\lambda^2}{128\pi^2} \frac{v_0^2}{\sqrt{e}} \equiv m_c^2 \tag{5.8}
\]
as the condition for condensation to be energetically favoured. At \( m^2 = m_c^2 \) the symmetric phase at \( v_B = 0 \) and the condensate phase at \( v_B^2 = v_0^2/\sqrt{e} \) have equal energy density, allowing the possibility of co-existence of the two phases.

The Higgs-boson mass \( M_h \) corresponds to \( M(\phi) \) in the physical vacuum, \( \phi_B = v_B \). Thus, from Eq. (4.20) we have

\[
M_h^2 \equiv M^2(\phi_B = v_B) = m^2 + \frac{1}{4}\lambda v_B^2. \tag{5.9}
\]

We shall neglect the \( m^2 \) term and justify this later. Noting that \( v_B^2 \) lies between \( e^{-1}v_0^2 \) and \( v_0^2 \), we see from (5.4) that

\[
M_h^2 = \mathcal{O}\left(\Lambda^2 \exp\left(-\frac{32\pi^2}{3\lambda}\right)\right). \tag{5.10}
\]

We want in the end to take the cutoff \( \Lambda \) to infinity, but such that the essential physics remains independent of \( \Lambda \). In this case the crucial condition is that the physical value of
the Higgs mass, \( M_h^2 \), should remain finite. The only way to obtain that result is to take \( \lambda \) to zero as \( \Lambda \to \infty \). Indeed, re-arranging (5.10) we see that \( \lambda \) must behave as:

\[
\frac{\lambda}{16\pi^2} \sim \frac{2}{3} \frac{1}{\ln(\Lambda^2/M_h^2)}.
\]

Thus, the coupling constant \( \lambda \) must depend on the cutoff, and, in particular, must tend to zero like \( 1/\ln \Lambda \). In that same limit \( v_B^2 \) must diverge as \( \ln \Lambda \) so that the product \( \lambda v_B^2 \), and hence \( M_h^2 \), is finite \[29\]. We also see now, from (5.8), that \( m^2 \leq m_c^2 \sim \lambda^2 v_B^2 \), so that \( m^2 = \mathcal{O}(\lambda M_h^2) \). Since \( \lambda \to 0 \), it follows that \( m^2 \) becomes vanishingly small, while \( M_h^2 \) remains finite. We were therefore justified in neglecting the \( m^2 \) term in Eq. (5.9). In summary, we have (for \( \Lambda \) in units of \( M_h \))

\[
\lambda = \mathcal{O}(1/\ln \Lambda), \quad m^2 = \mathcal{O}(1/\ln \Lambda), \quad v_B^2 = \mathcal{O}(\ln \Lambda).
\]

This type of behaviour ensures that the Higgs mass and the energy density at the minimum (5.7) are finite and so we have a physically significant limit when \( \Lambda \to \infty \). If \( \lambda \) were larger, then \( M_h \) would go to infinity (in particular, \( M_h = \mathcal{O}(\Lambda) \) for any small but finite value of \( \lambda \)). If \( m^2 \) were larger, then the term \( 4m^2\phi_B^2 \) would dominate \( V_{\text{eff}} \) which would have only a minimum at \( \phi_B = 0 \). The interesting region is close to the phase transition where the no-phion state at \( \phi_B = 0 \) and the condensate state at \( \phi_B = v_B \) are very close in energy.

In that region the elementary excitations of both vacua, phions and Higgs bosons, have vastly different masses as anticipated in the introduction.

Using the result (5.12) it is now straightforward to justify the neglect of the \( F(y) \) term in (4.25), at least for \( \phi_B \)'s that are comparable to \( v_B \), where it becomes \( \mathcal{O}(1/\ln \Lambda) \). At much smaller \( \phi_B \) the \( F(y) \) term does play some role; it serves to smooth out what would otherwise be a singularity in the fourth derivative at the origin.

\( V_{\text{eff}} \) has an important qualitative difference from the classical potential (1.1). The latter has a double-well form only for negative \( m^2 \) values and has a phase transition of second order at the value \( m^2 = 0 \). With \( V_{\text{eff}} \) the phase transition occurs at \( m^2 = m_c^2 \), Eq. (5.8), and the ‘order parameter’ \( \langle \Phi \rangle \) jumps from zero to \( e^{-1/4v_0} \). This first-order character, with the possibility of phase coexistence, remains in the limit \( \Lambda \to \infty \), despite the fact that \( m_c^2 \to 0 \) in units of \( M_h^2 \).

6. Phions and phonons

The effective potential in terms of the bare field \( \phi_B \) is an extremely flat function because the \( \phi_B = 0 \) vacuum and the \( \phi_B = \pm v_B \) vacua are infinitely far apart (\( v_B^2 \sim \ln \Lambda \) but
their energy densities differ only by a finite amount. To plot a graph of \( V_{\text{eff}}(\phi) \) one would naturally want to re-define the scale of the horizontal axis by defining a ‘renormalized’ or ‘re-scaled’ field \( \phi_R \). However, this does not correspond to a traditional wavefunction renormalization. Instead, it requires the following procedure \([11, 12]\): (i) Decompose the full field \( \Phi_B(x) \) into its zero- and finite-momentum pieces:

\[
\Phi_B(x) = \phi_B + h(x),
\]

where \( \int d^4x \ h(x) = 0 \). (ii) Re-scale the zero-momentum part of the field:

\[
\phi_B^2 = Z_\phi \phi_R^2
\]

with a \( Z_\phi \) that is large, of order \( \ln \Lambda \). (The condition determining \( Z_\phi \) is discussed below).

(iii) The finite-momentum modes \( h(x) \) remain unaffected.

This procedure is very natural in terms of the LHY calculation in Sect. 4, where singling out the \( k = 0 \) mode is crucial. From Eq. (4.5), excluding \( k = 0 \), we have

\[
h(x, t) = \sum_{k \neq 0} \frac{1}{\sqrt{2V_{E_k}}} \left[ a_k e^{ik \cdot x} + a_k^\dagger e^{-ik \cdot x} \right].
\]

Substituting the Bogoliubov transformation (4.13), re-organizing the summation using \( k \to -k \) in some terms, and finally using (4.16) we see that

\[
h(x, t) = \sum_{k \neq 0} \frac{1}{\sqrt{2V_{\tilde{E}_k}}} \left[ b_k e^{ik \cdot x} + b_k^\dagger e^{-ik \cdot x} \right].
\]

Thus, the finite-momentum part of the field is not re-scaled; it takes the canonical form in terms of phion or phonon variables. However, note that the Bogoliubov transformation (4.13) applies only to the \( k \neq 0 \) modes; “\( b_0, b_0^\dagger \)” remain undefined and are not necessarily related to \( a_0, a_0^\dagger \) (\( \sim \sqrt{N} \)) in the same way. Indeed, the \( a_k \)’s are discontinuous as \( k \to 0 \) because \( N \) phions occupy the \( k = 0 \) mode. However, for phonons the physics is continuous as \( k \to 0 \); that condition will fix \( Z_\phi \), as we show below.

In field-theory language the corresponding discussion is as follows. For the finite-momentum modes, general scattering-theory considerations lead to the Lehmann spectral decomposition (31), which implies that the wavefunction renormalization constant \( Z_h \) (in \( h_B(x) = \sqrt{Z_h} h_R(x) \)) must satisfy \( 0 < Z_h \leq 1 \), with \( Z_h \to 1 \) in the continuum limit if the theory is ‘trivial.’ However, these well-known arguments place no constraint on \( Z_\phi \), since there is no scattering theory for a zero-momentum mode — the incident particles would never reach each other. Instead, \( Z_\phi \) is fixed by the requirement that

\[
\left. \frac{d^2 V_{\text{eff}}}{d\phi_R^2} \right|_{\phi_R = v_R} = M_h^2.
\]

(6.5)
Although this is a familiar renormalization condition, the context here may be less familiar, so we would like to carefully explain its physical meaning in the ‘particle-gas’ language.

First, consider a slight perturbation of the symmetric vacuum state (“empty box”). We add a very small density \( n \) of phions, each with zero 3-momentum. The energy density is now:

\[
\mathcal{E}(n) = 0 + nm + \mathcal{O}(n^2 \ln n),
\]

where the first term is the energy of the unperturbed vacuum state (zero); the second term is the rest-mass cost of introducing \( N \) particles, divided by the volume; and the third term is negligible if we consider a sufficiently tiny density \( n \). Thus, we obviously have the relation

\[
\frac{\partial \mathcal{E}}{\partial n} \bigg|_{n=0} = m. \tag{6.7}
\]

The equivalent in field language follows from our previous relation

\[
n = \frac{1}{2} m \phi_B^2 \tag{6.8}
\]

and is given by

\[
\mathcal{E}(\phi_B) \equiv V_{\text{eff}}(\phi_B) = 0 + \frac{1}{2} m^2 \phi_B^2 + \mathcal{O}(\phi_B^4 \ln \phi_B^2),
\]

so that

\[
\frac{d^2 V_{\text{eff}}}{d\phi_B^2} \bigg|_{\phi_B=0} = m^2. \tag{6.10}
\]

Now, let us consider a slight perturbation of the broken-symmetry vacuum (the box filled with a spontaneously-generated condensate). Before we perturb it, this state has a density \( n_v \) of phions, where \( n_v \) is a (local) minimum of \( \mathcal{E}(n) \). From (6.8) we have the translation \( n_v = \frac{1}{2} m v_B^2 \). This vacuum state, though complicated in terms of phions, is simple in terms of the ‘phonon’ excitations corresponding to Higgs bosons: by definition it is just the state with no phonons. We now perturb it by adding a small density \( n' \) of phonons, each with negligibly small 3-momentum. (As noted above, phonons with zero momentum, created by \( b_0^\dagger \), are undefined; here we are effectively defining them by continuity.) The energy density of the perturbed state is then

\[
\mathcal{E}(n') = \mathcal{E}(n' = 0) + n'M_h + \ldots, \tag{6.11}
\]

where the first term is the energy density of the unperturbed state (\( = \mathcal{E}(n_v) \)); the second term is the rest-mass cost of the added phonons; and any other terms from phonon interactions are negligible if \( n' \) is small enough. Thus, paralleling (6.7) we have

\[
\frac{\partial \mathcal{E}}{\partial n'} \bigg|_{n'=0} = M_h. \tag{6.12}
\]
It is now natural to define a phonon field whose constant part, $f$, is related to the phonon density $n'$ by the analog of (6.8), namely

$$n' \equiv \frac{1}{2} M_h f^2.$$  
(6.13)

The “renormalized field” $\phi_R$ is simply this $f$ plus a constant. A constant must be added if we want to have $\phi_R$ proportional to $\phi_B$. Since, by definition, $f = 0$ when $\phi_B = v_B$, we need

$$\phi_R \equiv f + v_R$$  
(6.14)

with

$$\frac{v_R}{v_B} = \frac{\phi_R}{\phi_B} \equiv \frac{1}{Z \phi^{1/2}}.$$  
(6.15)

Now we can eliminate $f$ in favour of $\phi_R$ and re-write (6.13) as

$$n' = \frac{1}{2} M_h (\phi_R - v_R)^2.$$  
(6.16)

Hence, (6.11) can be re-written in field language as

$$\mathcal{E}(\phi_R) \equiv V_{\text{eff}}(\phi_R) = V_{\text{eff}}(\phi_R = v_R) + \frac{1}{2} M_R^2 (\phi_R - v_R)^2 + \ldots.$$  
(6.17)

The crucial condition, Eq. (6.5), follows directly from this. It just says that the phonon mass is, self-consistently, $M_h$. It is also, of course, the broken-vacuum counterpart of Eq. (6.10) for the symmetric vacuum.

The moral of this story is that the constant field $\phi_R - v_R$ is related to phonon density $n'$ in the same fashion that $\phi_B$ is related to the phion density $n$. Note that there is a duality under

$$\text{phions} \leftrightarrow \text{phonons}, \quad \phi_B \leftrightarrow \phi_R - v_R, \quad n \leftrightarrow n', \quad Z_\phi \leftrightarrow Z_\phi^{-1}.$$  
(6.18)

Physically, this means that, we may choose either phion or phonon degrees of freedom to describe the theory. Small excitations about the spontaneously broken vacuum are easily described in terms of phonons, but are complicated in terms of phions. Likewise, small excitations about the symmetric vacuum are easily described in terms of phions, but are complicated in terms of phonons.

### 7. Renormalized form of $V_{\text{eff}}$

Using the renormalized field $\phi_R = Z_\phi^{-1/2} \phi_B$ introduced in the last section we can write the effective potential in manifestly finite form. It is convenient to define a finite parameter $\zeta$
in terms of the physical mass and the renormalized vacuum expectation value:

\[ \zeta \equiv \frac{M^2_h}{8\pi^2 v_R^2}. \]  

(7.1)

Using \( M^2_h = \frac{4}{\lambda} \lambda v^2_R \) and \( Z_\phi = v^2_B/v^2_R \), one then has

\[ Z_\phi \equiv \frac{16\pi^2}{\lambda} \zeta, \]  

(7.2)

so that \( Z_\phi \) is of order \( 1/\lambda \), and hence of order \( \ln \Lambda \). Imposing the condition

\[ \left. \frac{d^2 V_{\text{eff}}}{d\phi^2_R} \right|_{\phi_R = v_R} = M^2_h, \]  

(7.3)

on the potential in Eq. (5.3) one finds

\[ m^2 = \frac{\lambda}{16\pi^2} \left( \frac{\zeta - 1}{2\zeta} \right) M^2_h, \]  

(7.4)

leading to the final form of \( V_{\text{eff}} \):

\[ V_{\text{eff}}(\phi_R) = \pi^2 \zeta(\zeta - 1)\phi_R^2 \left( 2v^2_R - \phi_R^2 \right) + \pi^2 \zeta^2 \phi^4_R \left( \ln \frac{\phi^2_R}{v^2_R} - \frac{1}{2} \right). \]  

(7.5)

The two independent quantities \( \zeta \) and \( v^2_R \) provide an intrinsic parametrization of the effective potential and replace the two bare parameters \((m^2, \lambda)\) of the original Hamiltonian. Both the mass and the vacuum energy density can be expressed in terms of these parameters:

\[ M^2_R = 8\pi^2 \zeta v^2_R, \]  

(7.6)

\[ V_{\text{eff}}(\phi_R = \pm v_R) = -\frac{\pi^2}{2} \zeta(2 - \zeta)v^4_R. \]  

(7.7)

The values \( \phi_R = \pm v_R \) are local minima of the effective potential for all positive values of \( \zeta \). However, only for \( \zeta < 2 \) do these non-trivial minima have lower energy density than the symmetric vacuum. Thus, the symmetry-breaking phase transition occurs at \( \zeta = 2 \) (corresponding to the value \( m^2 = m^2_c \)). At \( \zeta = 1 \) one reaches the massless case (or ‘Coleman-Weinberg regime’ \([\text{I}]\)) where \( m^2 = 0 \). In the range \( 2 \geq \zeta \geq 1 \), where spontaneous symmetry breaking happens even for a positive physical phion mass \( m \), the Higgs mass lies in the range

\[ 4\pi v_R \geq M_h \geq 2\sqrt{2\pi} v_R. \]  

(7.8)

Finally, the range \( 1 > \zeta > 0 \) corresponds to negative values of \( m^2 \) (‘tachionic phions’) where \( M^2_h \) can become arbitrarily small in units of \( v^2_R \). In the extreme case \( \zeta \to 0 \) one recovers the classical-potential results.
It is important to stress that the final, renormalized result for the effective potential, (7.5), has a more general validity than the specific bare expression Eq. (5.1). For instance, the Gaussian approximation generates a different bare expression, but leads to exactly the same renormalized result [6]. The point is that, once the bare result is re-expressed in terms of the renormalized field through the condition (7.3), the coupling $\lambda$ no longer appears. Thus, the final result is the same in any approximation related to (5.1) by a replacement of the nominal coupling constant $\lambda$ by some effective coupling $\tilde{\lambda}$ [6, 11, 14, 15]. Furthermore, any sort of effective coupling, $\tilde{\lambda}$, defined by summing some class of 4-point diagrams, is naturally of the same size as the original $\lambda$, if the latter has a size $O(1/\ln \Lambda)$. The reason is that in Feynman graphs each loop generates at most a $\ln \Lambda$ factor [31]. Thus, when each vertex has a factor $\lambda \sim O(1/\ln \Lambda)$, any 4-point diagram, of arbitrary complexity, is at most of order $1/\ln \Lambda$. That is, $\tilde{\lambda} = c\lambda$, where the finite number $c$ depends on precisely which class of diagrams have been taken into account. In the same manner one can show that any contribution to $V_{\text{eff}}$ not absorbed by the $\lambda \to \tilde{\lambda}$ replacement will be suppressed by one or more powers of $1/\ln \Lambda$ [32]. This counting argument accords with our argument that, if the theory is ‘trivial,’ it must be possible to reabsorb all interaction effects into suitable redefinitions of the classical energy density term, $(\lambda/4!)\phi^4_B \to (\tilde{\lambda}/4!)\phi^4_B$, and of the mass $M^2(\phi_B) = \frac{1}{2}\lambda\phi^2_B \to \frac{1}{2}\tilde{\lambda}\phi^2_B$, which governs the zero-point energy contribution from the free-field fluctuations of the shifted field (see Fig. 3).

8. Summary: a hierarchy of length scales

We have reconsidered the symmetry-breaking phase transition in $\lambda\Phi^4$ theory from a fresh perspective. The physics can be understood intuitively in terms of actual particles and their interactions. We have shown that the energy density $E(n)$, for a given particle density $n$, consists of $n$, $n^2$, and $n^2 \ln n$ terms arising from rest masses, short-range 2-body repulsions, and long-range 2-body attractions, respectively. The crucial $\ln n$ factor arises because the $-1/r^3$ attraction is so long range that it generates a logarithmic infrared divergence that is tamed only by screening from the background density. The translation $n = \frac{1}{2}m\phi^2$ converts $E(n)$ to the field-theoretic effective potential $V_{\text{eff}}(\phi)$.

The resulting form of $V_{\text{eff}}(\phi)$ — a sum of $\phi^2$, $\phi^4$, and $\phi^4 \ln \phi^2$ terms — is exactly what one should expect in a ‘trivial’ theory, as we argued in the introduction. Moreover, we self-consistently find that the theory is ‘trivial’ because the scattering length $a$ tends to zero in the continuum limit.

Despite ‘triviality’ — in fact, because of ‘triviality’ — a rich hierarchy of length scales
emerges. This hierarchy is summarized in the figure below in terms of the small parameter
\[ \epsilon \equiv 1 / \ln(\Lambda / M_h), \]
which tends to zero in the continuum limit, \( \Lambda \to \infty \).

\[
\begin{array}{cccccc}
& v_B^{-1} & & v_R^{-1} & \\
n_{v^{-1/3}} & a & d & \xi_h & \xi \\
(\Lambda^{-1}) & \left( \frac{1}{8\pi m} \right) & \left( M_h^{-1} \right) & (m^{-1}) \\
e^{-1/\epsilon} & e^{1/2} & e^{1/6} & 1 & e^{-1/2}
\end{array}
\]

Figure 4: Schematic representation of the length-scale hierarchy.

The ‘unit of length’ here is the correlation length of the broken phase, \( \xi_h \equiv 1/M_h \), the inverse of the Higgs mass. The phion Compton wavelength, \( \xi \equiv 1/m \), is much longer since, at or near the phase transition, the phion mass \( m \) becomes infinitesimal; \( m^2/M_h^2 \sim \epsilon \). At the same time, the phion-phion coupling becomes infinitesimal: \( \lambda \sim \epsilon \). The phion-phion cross section due to short-range repulsion is proportional to the square of the scattering length \( a(E) = \frac{1}{8\pi E} \). Even for the lowest phion energy, \( E = m \) the scattering length \( a = \frac{1}{8\pi m} \) vanishes like \( \epsilon^{1/2} \) in length units set by \( \xi_h \equiv 1/M_h \).

The phion density at the minimum of the effective potential \( \phi_B = v_B \) is \( n_v = \frac{1}{2} mv_B^2 \), which is very large, \( \mathcal{O}(\epsilon^{-1/2}) \), in physical units. Hence, the average spacing between two phions in the condensate, \( d \equiv n_v^{-1/3} \), is tiny compared with \( \xi_h \): \( \frac{d}{\xi_h} \sim \epsilon^{1/6} \). It is because there is such a high density of phions that their tiny interactions produce a finite effect on the energy density. Nevertheless, the phions in the condensate are very dilute because \( n_v a^3 \sim \epsilon \). In other words, the average spacing between phions is much, much larger than their interaction size: \( d/a \sim \epsilon^{-1/3} \).

Upon translating back to field language, a crucial element of this picture is the large re-scaling of the zero-momentum (spacetime constant) part of the field. The normalizations of \( \phi_B \) and \( \phi_R \) are set by the conditions
\[
\left. \frac{d^2 V_{\text{eff}}}{d\phi_B^2} \right|_{\phi_B=0} \equiv m^2, \quad \left. \frac{d^2 V_{\text{eff}}}{d\phi_R^2} \right|_{\phi_R=v_R} \equiv M_h^2. \tag{8.1}
\]
Since the theory is “nearly” a massless, free theory, \( V_{\text{eff}} \) is a very flat function of \( \phi_B \), and so the re-scaling factor \( Z_\phi \) in \( \phi_B^2 = Z_\phi \phi_R^2 \) is large, \( \mathcal{O}(1/\epsilon) \). Thus, the length scale \( v_B^{-1} \) is of order \( \epsilon^{1/2} \), comparable to the scattering length \( a \), while the length scale \( v_R^{-1} \) is finite, comparable to \( M_h^{-1} \).

21
The existence of a $Z_\phi$, distinct from the wavefunction renormalization, $Z_h$, of the finite-momentum modes, is now supported by some direct lattice evidence [33]. A Monte-Carlo simulation of the Ising limit of $\lambda\Phi^4$ theory was used to measure (i) the zero-momentum susceptibility $\chi$:

$$\chi^{-1} = \frac{d^2V_{\text{eff}}}{d\phi_B^2} \bigg|_{\phi_B=v_B} \equiv \frac{M_h^2}{Z_\phi},$$

(8.2)

and (ii) the propagator of the shifted field (at Euclidean momenta $p \neq 0$). The latter data was fitted to the form

$$G(p) = \frac{Z_h}{p^2 + M_h^2}$$

(8.3)

to obtain the mass and wavefunction-renormalization constant $Z_h$. The resulting $Z_h$ is slightly less than one, and seems to approach unity as the continuum limit is approached, consistent with the expected ‘triviality’ of the field $h(x)$ in the continuum limit. However, the $Z_\phi$ extracted from the susceptibility is clearly different. It shows a rapid increase above unity, and the trend is consistent with it diverging in the continuum limit.

This evidence, and the earlier lattice results for $V_{\text{eff}}$ [14, 17], provide objective support for our picture of symmetry breaking in $\lambda\Phi^4$ theory. As we have tried to show in this paper, the picture also has an appealing and very physical interpretation.

We would like to close with some speculations about some possible implications of these ideas in relation to gravity. Unlike other fields, which couple to gravity only via the $\sqrt{\det g}$ factor, a scalar field has a direct coupling to gravity through an $R\Phi^2$ term in the Lagrangian, where $R$ is the curvature scalar. For this reason, it has been proposed that Einstein gravity could emerge from spontaneous symmetry breaking [34, 35, 36]. The Newton constant, just like the Fermi constant, would then arise from the vacuum expectation value of a scalar field. It was suggested by van der Bij [37] that the scalar field inducing gravity could be the same scalar field responsible for electroweak symmetry breaking. The problem, though, is to understand the origin of the large re-scaling factor $\eta \sim 10^{34}$ needed in the coupling $\eta R \langle \Phi \rangle^2$ to obtain the Planck scale $\sim 10^{19}$ GeV from the Fermi scale $\sim 10^2$ GeV. Our results offer a possible solution to this puzzle. If we identify the Fermi scale with the physical vacuum field $v_R$, it is naturally infinitesimal with respect to the Planck scale, if we identify the latter with the bare condensate $v_B$. (This means that gravity must probe the scalar condensate at a much deeper level than electroweak interactions.) In this scenario we would need to back off from literally taking the $\Lambda \to \infty$ limit; instead we would need $\eta \equiv Z_\phi \sim \ln(\Lambda/M_h)$ to be large but finite, $\sim 10^{34}$. Note that this implies a cutoff $\Lambda$ that is enormously larger than the Planck mass [38].
One might wonder if gravity can still be neglected in discussing the phion dynamics. The formation of the phion condensate hinges on the very weak, long-range $-1/r^3$ potential, while gravity produces another weak — and even longer-range — interaction between phions. However, all is well provided that $Gm^2/r$ is much less than $A/r^3$ even for $r \sim r_{\text{max}}$. A little algebra shows that this condition is indeed satisfied, because the ratio $M_{\text{Planck}}^2/M_h^2$ (of order $\ln \Lambda$ in the above scenario) is much, much greater than unity.

A related issue is the ‘inflaton’ scalar field invoked in inflationary models of cosmology. The extraordinary fine-tuning of the scalar self-coupling needed to obtain a very slow rollover from the symmetric to the broken vacuum has led to the conclusion that “… the inflaton cannot be the Higgs field as had been originally hoped” \[^{40}\]. However, the $\phi_B$, $\phi_R$ distinction in our picture offers a natural way out of the difficulty. If gravity couples to the bare condensate, as postulated above, then it indeed sees an extremely flat effective potential. In our picture, a finite vacuum energy and a finite Higgs mass coexist with an infinitesimal slope of the effective potential (parametrized in terms of the bare vacuum field).

In our approach the natural order of magnitude relation is $M_h = O(v_R)$, with $m^2/M_h^2 \sim 1/Z_{\phi}$. In the scenario above, where $Z_{\phi} \sim 10^{34}$, we might expect a phion mass of order $m \sim 10^{-4} - 10^{-5}$ eV or smaller. Possibly, due to mixing effects with the graviton, this could produce deviations from the pure $1/r$ gravitational potential at the millimeter scale \[^{41}\] that can be tested in the next generation of precise ‘fifth-force’ experiments \[^{42}\].

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[10] Clearly, there is not a quadratic effective potential at the bare level; for one thing, the classical $\lambda \phi^4$ term must be present. One might postulate that after renormalization $V_{\text{eff}}$ — over the operative range of $\phi$ — becomes proportional to $\phi^2$ in the symmetric phase and to $(\phi \pm v)^2$ in the broken phase. However, then the question arises: What lies “in between” these results? How do they match up at the phase transition? After all, these are just two phases of the same theory. A broken-phase $V_{\text{eff}}$ proportional to $(\phi \pm v)^2$ would imply that no temperature, however high, could ever restore the symmetry; in that case one would have explicit, rather than spontaneous, symmetry breaking. We confess that the misconception that ‘triviality’ would imply a quadratic effective potential is the reason that, until 1993, we mistakenly thought that Gaussian-effective-potential results were in conflict with ‘triviality.’

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The point is that the virtual phions being exchanged propagate through a background medium and experience multiple collisions with background $k = 0$ phions. Each such collision corresponds to a “mass insertion” $-i\chi$ in the propagator, with $\chi = \lambda n/m = \frac{1}{2} \lambda \phi^2$. The resulting geometric sum changes the inverse propagator into $p^2 - M^2(\phi)$ with $M^2(\phi) = m^2 + \chi$. This conclusion is confirmed by the quasiparticle spectrum obtained in section 4.

More generally we mean this notation as a shorthand indication that any necessary mass-renormalization has been done; that is, a $\Phi^2$ and a $\Phi$-independent term are subtracted from $H$ such that the symmetric vacuum has zero energy, $\langle 0 | H | 0 \rangle = 0$, and such that $m$ corresponds to the physical mass of the phion particle.

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Fig. 1. The lowest order Feynman diagram of $\lambda \Phi^4$ theory corresponding to a $\delta^{(3)}(r)$ potential.

Fig. 2. The basic one-loop diagram describing phonon-phonon scattering to order $\lambda^2$. There are three distinct diagrams corresponding to the three Mandelstam variables $s$, $t$ and $u$. The long-range attractive $-1/r^3$ potential is generated by the Fourier transform of the $t$- and $u$-channel contributions.
Fig. 3. A pictorial representation of the structure of $V_{\text{triv}}$. This involves a simultaneous replacement of the bare coupling constant $\lambda$ in the classical energy density and in the zero-point energy of the shifted $h$-field. Dashed lines represent the $\phi_B$ constant vacuum field and solid lines the $h$-field.