Crucial Dependence of “Precarious” and “Autonomous” $\phi^4$'s
Upon the Normal-ordering Mass

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Using the Gaussian wave-functional approach with the normal-ordering renormalization prescription, we show that for the (3+1)-dimensional massive $\lambda\phi^4$ theory, “precarious” and “autonomous” $\phi^4$'s can exist if and only if the normal-ordering mass is equal to the classical masses at the symmetric and asymmetric vacua, respectively.

The triviality problem of quantum field theories, for example, the $\lambda\phi^4$ theories, is very important and has been receiving lots of investigations since the early 1970’s [1]. In general, it is believed that $(D+1)$-dimensional field theories $(D \geq 3)$ have been proved to be trivial (i.e., to be non-interacting or else inconsistent). In the 1980’s, nevertheless, P. M. Stevenson and his collaborators took a new view of the $(3+1)$-dimensional $\lambda\phi^4$ theory and accordingly proposed two non-trivial $\lambda\phi^4$ theories: “percarious $\phi^4$” and “autonomous $\phi^4$”. Within the framework of the Gaussian wave-functional approach (GWFA), for the case of the negative and infinitesimal bare coupling, Stevenson renormalized the $(3+1)$-dimensional $\lambda\phi^4$ theory by a set of renormalization conditions of the mass and coupling parameters at the zero vacuum expectation value (VEV) of field operator, and obtained a non-trivial theory called “precarious $\phi^4$” [2]: for the case of the positive and infinitesimal bare coupling, Stevenson and Tarrach also renormalized the $(3+1)$-dimensional $\lambda\phi^4$ theory by another set of renormalization conditions of the mass and coupling parameters at the zero VEV of field operator and further by an infinite wavefunction renormalization condition, and obtained the other non-trivial theory called “autonomous $\phi^4$” [3]. “Precarious $\phi^4$” has a stable, symmetric phase [4], whereas “autonomous” $\phi^4$ can exhibit either a massless symmetry phase or a massive broken-symmetry phase [5]. Since the two non-trivial theories were proposed, they have attracted much attention. These two non-trivial theories were demonstrated to exist in O(N)-symmetric $\lambda\phi^4$ theory [6], and their generalizations to many other complicated models were made, too [7]. Furthermore, it was shown that in the context of dimensional regularization, “percarious $\phi^4$” arises as the $(D+1)\rightarrow (4)_+\text{ limit of the } (D+1)$-dimensional $\lambda\phi^4$ theory with $D > 3$, whereas “autonomous $\phi^4$” can be understood as the $(D+1)\rightarrow (4)_-\text{ limit of the } (D+1)$-dimensional $\lambda\phi^4$ theory with $D < 3$ [8]. Again, “precarious” and “autonomous” $\phi^4$'s were discussed by some methods beyond the Gaussian approximation, and the existence of “autonomous $\phi^4$” was still demonstrated but those methods did not all confirm the appearance of “precarious $\phi^4$” [9]. Besides, “autonomous $\phi^4$” arising from the classical pure $\lambda\phi^4$ theory (massless) was believed to not conflict with “Triviality”, and used in relevant Models to predict a 2.2 Tev Higgs boson [10].

To our knowledge, no work discussed the dependences of “precarious” and “autonomous” $\phi^4$’s upon renormalization points. This is perhaps because results obtained from the GWFA with the standard renormalization prescription [11] are exactly renormalization-group invariant [12] (1985, Sect.III). However, although “precarious $\phi^4n$” was obtained according to the standard renormalization prescription, an approximation was made in renormalizing the coupling parameter [13], and as for the “autonomous $\phi^4n$”, an unusual set of renormalization conditions was proposed within the framework of the GWFA [14]. On the other hand, as was mentioned above, it was at a peculiar renormalization point, the zero VEV of field operator, that both “precarious” and “autonomous” $\phi^4$’s were demonstrated to exist [15]. Therefore, it is necessary and important to consider whether “precarious” and “autonomous” $\phi^4$’s would exist or not when the $(3+1)$-dimensional $\lambda\phi^4$ theory is renormalized at any other renormalization points. In this letter, we intend to address this problem.

For the above problem, the renormalization prescriptions utilized by Stevenson et al [16] are not too convenient because it must be executed at a explicit renormalization point given previously. Nevertheless, within the framework of the GWFA, executing the standard renormalization prescription [11] is equivalent to executing Coleman’s renormalization prescription [10] [13], and further, executing a renormalization procedure at different renormalized points amounts to choosing different values of the normal-ordering mass in Coleman’s normal-ordering renormalization prescription [11]. Thus, if Coleman’s normal-ordering renormalization prescription is adopted, the dependence of a field theory on renormalization points will be equivalent to the dependence of that theory on the normal-ordering mass. Luckily, Coleman’s renormalization prescription can be performed at any normal-ordering mass which needn’t be explicitly

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1In the standard renormalization prescription of the GWFA, the definitions of renormalized physical parameters resemble those in the classical theory concerned.
given, and this is evidently convenient for the purpose here. Hence, in order to know if the existences of “precarious” and “autonomous” φ's depend upon renormalization points, it is sufficient to renormalize the (3+1)-dimensional λφ⁴ theory on the basis of Coleman's normal-ordering renormalization prescription and pay attention to the dependence of “precarious” and “autonomous” φ's upon values of the normal-ordering mass instead of the renormalization points.

In the following, based on Coleman's normal-ordering renormalization prescription, we shall analyse the Gaussian effective potential (GEP) of the (3+1)-dimensional massive λφ⁴ theory, manage to turn the GEP finite, and accordingly show that “precarious” and “autonomous” φ's appear only at two particular values of the normal-ordering mass, respectively.

Consider the massive single-component λφ⁴ field theory in (3+1) dimensions,  
\[ L = \frac{1}{2} \partial_{\mu} \phi_{x} \partial^{\mu} \phi_{x} - \frac{1}{2} m^{2} \phi_{x}^{2} - \lambda \phi_{x}^{4}, \tag{1} \]
where ϕₓ ≡ φ(⃗x) and ⃗x is the position vector in the three-dimensional space. In the fixed-time functional Schrödinger picture, the Hamiltonian reads
\[ H = \int_{x} H_{x} = \int_{x} \left\{ \frac{1}{2} \Pi_{x}^{2} + \frac{1}{2} (\nabla \phi_{x})^{2} + \frac{1}{2} m^{2} \phi_{x}^{2} + \lambda \phi_{x}^{4} \right\}, \tag{2} \]
where, \( f_{x} = \int d^{3} \bar{x}, \nabla \) is the gradient operator and \( \Pi_{x} = -i \frac{\delta}{\delta \phi_{x}} \) conjugate to the field operator \( \phi_{x} \). The mass and coupling parameters \( m \) and \( \lambda \) govern the classical physics of the system. When \( m^{2} > 0 \) the classical vacuum is symmetric and the classical mass squared is \( m^{2} \), but when \( m^{2} < 0 \) the classical vacuum is spontaneously symmetry-breaking and the classical mass squared is \(-2m^{2}\).

Take as an ansatz the general Gaussian wave-functional  
\[ |\varphi > \rightarrow \Psi[\phi; \varphi, P, f] = N_{f} \exp \{ i \int P_{x} \phi_{x} - \frac{1}{2} \int_{x,y} (\phi_{x} - \varphi_{x}) f_{xy}(\phi_{y} - \varphi_{y}) \}, \tag{3} \]
where \( P_{x}, \varphi_{x} \) and \( f_{xy} \) are variational parameter-functions. \( N_{f} \) is some normalization constant, and depends upon \( f_{xy} \). With respect to any normal-ordering mass \( M \) (an arbitrary positive constant with the dimension of mass), one can Nomal-order the Hamiltonian density in Eq.(2) according to Ref. [11], and has
\[ \mathcal{N}_{M}[H_{x}] = H_{x} - 3\lambda I_{1}(M^{2})\phi_{x}^{2} + \frac{3}{4} \lambda I_{2}^{2}(M^{2}) - \frac{1}{4} m^{2} I_{1}(M^{2}) - \frac{1}{2} I_{0}(M^{2}) + \frac{1}{4} M^{2} I_{1}(M^{2}) \tag{4} \]
with the notation 
\[ I_{n}(M^{2}) = \int \frac{d^{3} p}{(2\pi)^{3}} \sqrt{p^{2} + M^{2}}. \tag{5} \]
Here, \( p = |\vec{p}| \), and \( \mathcal{N}_{M}[\cdots] \) denotes normal-ordering with respect to \( M \). Following the Refs. [11] [13], one can first calculate the energy \( \int_{x} < \varphi | \mathcal{N}_{M}[H_{x}] | \varphi > \), then take \( \varphi_{x} \) as a constant \( \varphi \), and finally, minimize variationally the energy with respect to \( P \) as well as \( f \) to obtain the GEP. Consequently, \( P_{x} = 0 \), the Fourier component of \( f_{xy} \) is \( f(p) = \sqrt{p^{2} + 4\lambda^{2} \varphi^{2}} \), and the GEP has the following expression
\[ V(\varphi) = \frac{1}{2} [I_{0}(\Omega^{2}) - I_{0}(M^{2})] - \frac{1}{4} [\Omega^{2} I_{1}(\Omega^{2}) - M^{2} I_{1}(M^{2})] \]
\[ + \frac{1}{4} m^{2} [I_{1}(\Omega^{2}) - I_{1}(M^{2})] + \frac{3}{4} \lambda [I_{0}(\Omega^{2}) - I_{1}(M^{2})]^{2} \]
\[ + 3\lambda [I_{1}(\Omega^{2}) - I_{1}(M^{2})] \varphi^{2} + \frac{1}{2} m^{2} \varphi^{2} + \lambda \varphi^{4} \tag{6} \]
with the gap \( \Omega = \Omega(\varphi) \) satisfying
\[ \Omega^{2} = m^{2} + 6\lambda [I_{1}(\Omega^{2}) - I_{1}(M^{2})] + 12\lambda \varphi^{2}. \tag{7} \]

For the lower-dimensional cases (\( D < 3 \)), the corresponding GEP contains no divergences [13], and no further renormalization procedure need to be considered. However, in the present case, the situation is completely different and the right hands of Eqs.(6) and (7) are still full of divergences. Hence, we have to find a further renormalization scheme for making the GEP finite. Next, let us analyse those divergences appeared in Eqs.(6) and (7).
A straightforward calculation can yield
\[
I_1(\Omega^2) - I_1(M^2) = -\frac{1}{8\pi^2}(\Omega^2 - M^2) + \frac{1}{8\pi^2}\Omega^2\ln\frac{\Omega^2}{M^2} - \frac{1}{2}(\Omega^2 - M^2)I_2(M^2)
\]
(8)
and
\[
\frac{1}{2}[I_0(\Omega^2) - I_0(M^2)] - \frac{1}{4}(\Omega^2)I_1(\Omega^2) - M^2I_1(M^2)]
= \frac{1}{128\pi^2}(\Omega^4 - M^4) - \frac{1}{64\pi^2}\Omega^2\ln\frac{\Omega^2}{M^2} + \frac{1}{16}(\Omega^4 - M^4)I_2(M^2),
\]
(9)
where \(I_2(M^2)\) is a logarithmic divergent integral. Substituting the last two equations into Eqs.(6) and (7), one can see that a logarithmic divergence exists in Eq.(7) and the right hand of Eq.(6) has the terms with logarithmic divergences squared at most, except for divergent constants (independent of \(\varphi\)). It is well known that in the case of lower dimensions, Coleman’s normal-ordering renormalization prescription amounts to the renormalization of the mass parameter \(I\) (1985), and the relation between the bare and renormalized coupling parameters is in fact a finite relation \(I\) (1985). Thus, generally, for the present case, in order to get rid of the divergences in Eqs.(6) and (7), perhaps we should consider further a real renormalization of the coupling parameter and an infinite renormalization or infinitely re-scaled that those divergences both in Eq.(6) and in Eq.(7) can cancel out, respectively, except for divergent constants in Eq.(6).

From the above analysis, we take
\[
\lambda = a_1I_2^{-1}(M^2) + a_2I_2^{-2}(M^2), \quad \varphi^2 = bI_2(M^2)\Phi^2
\]
(10)
with \(a_1\) and \(b\) being some constants to be determined. Note that for \(b\), what is necessary is only to determine it up to a finite rescaling of \(\Phi\). In the ansatz Eq.(10), \(a_2\) should be a parameter relevant to renormalized counterpart of the bare coupling \(\lambda\), with a finite relation, if Eq.(10) can really renormalize the theory. Analysing Eqs.(6) and (7), one can see that the ansatz Eq.(10) should has contained all possibilities to remove various divergences both in Eqs.(6) and (7), and other ansatzes would not either remove those divergences or lead to a new result. Substituting Eq.(10) into Eq.(6), one can find that those divergences in Eq.(7) cancel out, and Eq.(6) has only two types of divergent terms: \(\Phi^2I_2(M^2)\) and \(\Phi^2I_2(M)\), except for the divergent constant \(D \equiv V(\varphi = 0)\). Then taking the coefficients of \(\Phi^2I_2(M)\) and \(\Phi^4I_2(M)\) terms as zero, respectively, leads to
\[
\begin{cases}
\frac{b(m^2 + 3a_1)M^2}{4(1+3a_1)} = 0 \\
\frac{a_1b^2}{1+3a_1}(1 - 6a_1) = 0
\end{cases}
\]
(11)
and consequently the GEP in Eq.(6) also contains no divergences, except for \(D\). Thus, the solutions for the set of equations Eq.(11) will be helpful to determine the renormalization schemes.

Obviously, there are three solutions for the set of equations Eq.(11):

\[(i). \quad a_1 \sim I_2(M^2), \quad b = I_2^{-1}(M^2); \]
\[(ii). \quad a_1 = -\frac{1}{3}, \quad b = I_2^{-1}(M^2), \quad M^2 = m^2; \]
\[(iii). \quad a_1 = \frac{1}{6}, \quad M^2 = -2m^2, \quad b \text{ is finite}. \]
(12)
In the strict sense of mathematics, \(ii\) in Eq.(12) cannot be regarded as a solution of Eq.(11), because in the second equation of Eq.(11) it leads to the infinitesimal \(b^2\) dividing by zero \(3\). Nevertheless, \(ii\) in Eq.(12) can really get rid of those divergences in Eqs.(6) and (7). In fact, in Eqs.(6) and (7) with Eq.(10), there is not the denominator \((1+3a_1)\) at all, and the appearance of this denominator is due to a mathematical deformation of rewriting Eq.(7). Additionally, although the solution \(i\) can make Eqs.(6) and (7) have no explicit divergences (of course except for a divergent constant in Eq.(6)), it means that \(\lambda\) is finite and therefore does not produce a viable theory, for the vacuum could become unstable \(\mathbb{1}\). In view of this, in the following, we discuss the other two solutions.

The solution \(ii\) in Eq.(12) implies that it is not necessary to perform the wave-function renormalization. For the convenience of comparison with Ref. \(\mathbb{1}\), let \(a_2 = \frac{k}{12\pi}\) with \(k = \frac{4\pi^2}{\lambda_R}\). The renomalization scheme is

\(^{2}\)The author thanks the referee for commenting upon this point.
and

\[ \lambda = \frac{1}{3} I_2^{-1}(M^2) + \frac{1}{12\pi^2} I_2^{-2}(M^2) . \] (13)

Substituting Eqs.(13), (8) and (9) into Eqs.(6) and (7), we have

\[ V(\Phi) = \frac{1}{2} \Omega^2 \Phi^2 + \frac{1}{128\pi^2} [2\Omega^4 \ln \frac{\Omega^2}{m^2} - (3 - 2k)(\Omega^2 - m^2)^2] + D_1 \] (14)

and

\[ \Omega^2 \ln \frac{\Omega^2}{m^2} = (\Omega^2 - m^2)(1 - k) - 16\pi^2 \Phi^2 \] (15)

with \( D_1 \) being a divergent constant. Eqs.(14) and (15) are identical to Eqs.(5.9) and (5.10) in Ref. [3](1985), respectively. Thus, we reproduce "precarious \( \phi^4 \)". However, as has been seen in the above, only when \( M^2 = m^2 \), can the divergences be removed, and therefore, if and only if \( M^2 = m^2 \), "precarious \( \phi^4 \)" exists. By the way, because \( M^2 \) must be positive, "precarious \( \phi^4 \)" corresponds to Eq.(1) with \( m^2 > 0 \). That is to say, if and only if the classical mass in the case of \( m^2 > 0 \) and \( \lambda < 0 \) is taken as the normal-ordering mass, "precarious \( \phi^4 \)" can appear. This also implies that "precarious \( \phi^4 \)" arises from only the symmetrical phase of the classical theory.

Now we are in a position to discuss the solution (iii). In this case, we take \( \alpha_2 = -\frac{1}{6} \lambda_R \) and \( b = \frac{1}{2} \) for the consistence with Ref. [3]. Hence, we have the renormalization scheme

\[ \begin{cases} 
\lambda = \frac{1}{6} I_2^{-1}(M^2) - \frac{1}{F} \lambda_R I_2^{-2}(M^2) \\
\phi^2 = \frac{1}{F} I_2(M^2) \phi^2 \\
M^2 = -2m^2 
\end{cases} \] (16)

The third equation in Eq.(16) requires that \( m^2 \) is negative. Employing this scheme, Eqs.(6) and (7) can be expressed as

\[ V(\Phi) = \left( \frac{1}{6} \lambda_R - \frac{1}{24\pi^2} \right) m^2 \Phi^2 + \frac{1}{36} \lambda_R \Phi^4 + \frac{1}{144\pi^2} \Phi^4 \left( \ln \frac{\Omega^2}{2m^2} - \frac{3}{2} \right) + D_2 \] (17)

and

\[ \Omega^2 = \frac{2}{3} \Phi^2 \] (18)

with \( D_2 \) the divergent constant. Eq.(17) is the GEP obtained by the scheme Eq.(16). A numerical analysis indicates that so long as \( \lambda_R \) is not too small the SSB remains to exist, whereas if \( \lambda_R \) is small the vacuum enjoys its symmetry 3.

Owing to \( I_2(x) - I_2(y) = -\frac{1}{3} \ln \frac{y^2}{x^2} \), we can introduce a finite characteristic scale \( \mu \) with the dimension of mass, which is related to the renormalized coupling with \( \lambda_R = -\frac{1}{4\pi^2} \ln \frac{\mu^2}{m^2} \). Rewriting \( \lambda \) in Eq.(16) in terms of \( \mu \):

\[ \lambda = \frac{1}{6} \frac{m^2}{I_2(\mu^2)} \] (19)

we obtain

\[ V(\Phi) = \frac{M^2}{48\pi^2} \left( 1 + \ln \frac{\mu^2}{M^2} \right) \Phi^2 + \frac{1}{144\pi^2} \Phi^4 \left( \ln \frac{\Omega^2}{\mu^2} - \frac{3}{2} \right) . \] (20)

It is evident that the expression of Eq.(20) is identical to that of "autonomous \( \phi^4 \)" in Ref. [3] when setting \( m_0^2 = \frac{\Omega^2}{24\pi^2} \). Here, we have seen that in the present case, if and only if we make a special choice of normal-ordering mass, i.e., \( M^2 = -2m^2 \), the GEP is renormalized and "autonomous \( \phi^4 \)" exist. Contrary to "precarious \( \phi^4 \)", "autonomous \( \phi^4 \)" corresponds to Eq.(1) with \( m^2 < 0 \), and originates from the asymmetrical phase of the classical theory. That is to say, only when the classical mass in the case of \( m^2 < 0 \) and \( \lambda > 0 \) is chosen as the normal-ordering mass, "autonomous \( \phi^4 \)" appears.

Perhaps one must have noticed that \( \lambda \) in Eq.(16) or Eq.(19) is just Eq.(9) in Ref. [3] 4. In the same way as Eq.(19) is derived, one also can rewrite \( \lambda \) in Eq.(13) and find that \( \lambda \) in Eq.(13) is just Eq.(5.5) in Ref. [3](1985). Further, simple

\[ ^3 \text{Note that } I_2(\mu^2) \text{ here is equal to } 2I_{1-1}(\mu) \text{ in Ref. [3].} \]
calculations of $\Omega(\varphi = 0)$ from Eq.(7) and $\lambda$ in Eqs.(13) and (16) indicate that $M^2 = m^2$ is really consistent with the renormalization condition of the mass in Ref. 1, and $M^2 = -2m^2$ consistent with Eq.(10) of Ref. 3. So the renormalization scheme Eq.(13) is equivalent to that in Ref. 3, and the renormalization scheme Eq.(16) equivalent to the scheme Eqs.(9),(10) and (11) in Ref. 3. That is to say, the two values of the normal-ordering mass in the treatment here, $m$ for the case of $m^2 > 0$ and $\lambda < 0$ and $\sqrt{-2m^2}$ for the case of $m^2 < 0$ and $\lambda > 0$, correspond to the renormalization point $\varphi = 0$ chosen for the case of $\lambda < 0$ in Ref. 3 and for the case of $\lambda > 0$ in Ref. 3, respectively. Because other choices of the normal-ordering mass can not yield a renormalized and non-trivial theory, the existence of both “precarious $\phi^4$” and “autonomous $\phi^4$” is crucially dependent upon the normal-ordering mass (or renormalization points).

In the same way, we have also considered the O(N)-symmetric $\lambda\phi^4$ theory and obtained an analogous conclusion. That is, for the O(N)-symmetric $\lambda\phi^4$ theory, “precarious” and “autonomous” $\phi^4$s in Ref. 4 exist only at $M^2 = m^2 (m^2 > 0)$ and at $M^2 = -\frac{2m^2}{-1+N/3} (m^2 < 0)$, respectively. Note that the value $M^2 = m^2$ is just the classical mass with the symmetric vacuum and $M^2 = -\frac{2m^2}{-1+N/3}$ the one with the asymmetric vacuum. Thus, we conclude that the existence of “precarious” and “autonomous” $\phi^4$s is uniquely at the respective particular values of the normal-ordering mass the respective classical masses. Besides, we want to emphasize that as has been shown in the derivation of this paper, if any other positive value (not the classical mass) is chosen as the normal-ordering mass, not only the “precarious” and “autonomous” $\phi^4$s disappear, but also the massive $\lambda\phi^4$ with (3+1) dimensions is either trivial or non-renormalizable for any $\lambda$ within the framework of the GWFA. Finally, it should be mentioned that because normal-ordering mass must be greater than zero, the conclusion in this paper is not valid for the massless $\lambda\phi^4$ theory.

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\[4\]In the case of $m^2 < 0$, $I_1(\Omega)$ in $\Omega(\varphi = 0)$ should be taken as $I_1(0)$ to keep consistence with Ref. 3.
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