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Research Article

The Higgs Boson: From the Lattice to LHC

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We discuss the triviality and spontaneous symmetry breaking scenario where the Higgs boson without self-interaction coexists with spontaneous symmetry breaking. We argue that nonperturbative lattice investigations support this scenario. Moreover, from lattice simulations, we predict that the Higgs boson is rather heavy. We estimate the Higgs boson mass $m_H = 754 \pm 20$ (stat) ± 20 (syst) GeV and the Higgs total width $\Gamma(H) \simeq 340\,\mathrm{GeV}$.

1. Introduction

A cornerstone of the Standard Model is the mechanism of spontaneous symmetry breaking that, as is well known, is mediated by the Higgs boson. Then, the discovery of the Higgs boson is the highest priority of the Large Hadron Collider (LHC) [1, 2].

Usually the spontaneous symmetry breaking in the Standard Model is implemented within the perturbation theory which leads to predict that the Higgs boson mass squared, m_H^2, is proportional to $\lambda_R v_R^2$, where v_R is the known weak scale (246 GeV) and λ_R is the renormalized scalar self-coupling. However, it has been conjectured since long time [3] that self-interacting four dimensional scalar field theories are trivial, namely, $\lambda_R \to 0$ when $\Lambda \to \infty$ (Λ ultraviolet cutoff). Even though no rigorous proof of triviality exists, there exist several results which leave little doubt on the triviality conjecture [4–7]. As a consequence, within the perturbative approach, these theories represent just an effective description, valid only up to some cut-off scale Λ, for without a cutoff, there would be no scalar self-interactions and without them no symmetry breaking. However, within the variational Gaussian approximation, it has been suggested in [8] that this conclusion could not be true. The point is that the Higgs condensate and its quantum fluctuations could undergo different rescalings when changing the ultraviolet cutoff. Therefore, the relation between m_H and the physical v_R is not the same as in perturbation theory. Indeed, according to this picture, one expects that the condensate rescales as $Z_\varphi \sim \ln\Lambda$ in such a way to compensate the $1/\ln\Lambda$ from λ_R. As a consequence, the ratio m_H/v_R would be a cutoff-independent constant. In other words, one should have

$$m_H = \xi v_R, \tag{1.1}$$

where ξ is a cutoff-independent constant.

It is noteworthy to point out that (1.1) can be checked by nonperturbative numerical simulations of self-interacting four dimensional scalar field theories on the lattice. Indeed, in previous studies [9, 10], we found numerical evidences in support of (1.1). Moreover, our numerical results showed that the extrapolation to the continuum limit leads to the quite simple result:

$$m_H \simeq \pi v_R, \tag{1.2}$$

pointing to a rather massive Higgs boson without self-interactions (triviality).

The plan of the paper is as follows. In Section 2, we illustrate that triviality could coexist with spontaneous symmetry breaking within the simplest self-interacting scalar field theory in four dimensions. In Section 3, we briefly review the lattice indications for the nonperturbative interpretation of triviality in self-interacting four-dimensional scalar field theories and furnish our best numerical determination of the constant ξ in (1.1). Section 4 is devoted to discuss some experimental signatures of the Higgs boson at LHC. Finally, our conclusions are drawn in Section 5.

2. Triviality and Spontaneous Symmetry Breaking

In this section, we discuss the triviality and spontaneous symmetry breaking scenario within the simplest scalar field theory, namely, a massless real scalar field Φ with quartic self-interaction $\lambda\Phi^4$ in four dimensions:

$$\mathcal{L} = \frac{1}{2}\left(\partial_\mu \Phi_0\right)^2 - \frac{1}{4}\lambda_0 \Phi_0^4, \tag{2.1}$$

where λ_0 and Φ_0 are the bare coupling and field, respectively. As it is well known [11, 12], the one-loop effective potential is given by summing the vacuum diagrams:

$$V_{1-\text{loop}}(\phi_0) = \frac{1}{4}\lambda_0\phi_0^4 - \frac{i}{2}\int \frac{d^4k}{(2\pi)^4} \ln\left[-k_0^2 + \vec{k}^2 + 3\lambda_0\phi_0^2 - i\epsilon\right]. \tag{2.2}$$

Integrating over k_0 and discarding a (infinite) constant give

$$V_{1-\text{loop}}(\phi_0) = \frac{1}{4}\lambda_0\phi_0^4 + \frac{1}{2}\int \frac{d^3k}{(2\pi)^3}\sqrt{\vec{k}^2 + 3\lambda_0\phi_0^2}. \tag{2.3}$$

This last equation can be interpreted as the vacuum energy of the shifted field:

$$\Phi_0 = \phi_0 + \eta \tag{2.4}$$

in the quadratic approximation. Indeed, in this approximation, the hamiltonian of the fluctuation η over the background ϕ_0 is

$$\mathscr{H}_0 = \frac{1}{2}(\Pi_\eta)^2 + \frac{1}{2}\left(\vec{\nabla}\eta\right)^2 + \frac{1}{2}\left(3\lambda_0\phi_0^2\right)\eta^2 + \frac{1}{4}\lambda_0\phi_0^4. \tag{2.5}$$

Introducing an ultraviolet cutoff Λ, we obtain, from(2.3)

$$V_{1-\text{loop}}(\phi_0) = \frac{1}{4}\lambda_0\phi_0^4 + \frac{\omega^4}{64\pi^2}\ln\left(\frac{\omega^2}{\Lambda^2}\right), \qquad \omega^2 = 3\lambda_0\phi_0^2. \tag{2.6}$$

It is easy to see that the one-loop effective potential displays a minimum at

$$3\lambda_0 v_0^2 = \frac{\Lambda^2}{\sqrt{e}}\exp\left[-\frac{16\pi^2}{9\lambda_0}\right]. \tag{2.7}$$

Moreover,

$$V_{1-\text{loop}}(v_0) = -\frac{\omega^4}{128\pi^2}, \tag{2.8}$$

so that

$$V_{1-\text{loop}}(\phi_0) = \frac{\omega^4}{64\pi^2}\left[\ln\left(\frac{\phi_0^2}{v_0^2}\right) - \frac{1}{2}\right]. \tag{2.9}$$

According to the renormalization group invariance, we impose that, for $\Lambda \to \infty$,

$$\left[\Lambda\frac{\partial}{\partial\Lambda} + \beta\frac{\partial}{\partial\lambda_0} + \gamma\phi_0\frac{\partial}{\partial\phi_0}\right]V_{1-\text{loop}}(\phi_0) = 0. \tag{2.10}$$

Within perturbation theory, one finds

$$\gamma = 0, \qquad \beta = \frac{9}{8\pi^2}\lambda_0^2. \tag{2.11}$$

Thus, the one-loop corrections have generated spontaneous symmetry breaking. However, the minimum of the effective potential lies outside the expected range of validity of the one-loop approximation, and it must be rejected as an artefact of the approximation. On the other hand, as discussed in Section 1, there is no doubt on the triviality of the theory. As a consequence, within perturbation theory, there is no room for symmetry breaking. However, following the suggestion of [8], we argue below that spontaneous symmetry breaking could be compatible with triviality. The arguments go as follows. Write

$$\Phi_0 = \phi_0 + \eta, \tag{2.12}$$

where ϕ_0 is the bare uniform scalar condensate; thus, triviality implies that the fluctuation field η is a free field with mass $\omega(\phi_0)$. This means that the *exact* effective potential is

$$V_{\text{eff}}(\phi_0) = \frac{1}{4}\lambda_0\phi_0^4 + \frac{1}{2}\int \frac{d^3k}{(2\pi)^3}\sqrt{\vec{k}^2 + \omega^2(\phi_0)} = \frac{1}{4}\lambda_0\phi_0^4 + \frac{\omega^4(\phi_0)}{64\pi^2}\ln\left(\frac{\omega^2(\phi_0)}{\Lambda^2}\right). \tag{2.13}$$

Moreover, the mechanism of spontaneous symmetry breaking implies that the mass of the fluctuation is related to the scalar condensate as

$$\omega^2(\phi_0) = 3\tilde{\lambda}\phi_0^2, \qquad \tilde{\lambda} = a_1\lambda_0, \tag{2.14}$$

where a_1 is some numerical constant.

Now, the problem is to see if it exists the continuum limit $\Lambda \to \infty$. Obviously, we must have

$$\left[\Lambda\frac{\partial}{\partial\Lambda} + \beta(\lambda_0)\frac{\partial}{\partial\lambda_0} + \gamma(\lambda_0)\phi_0\frac{\partial}{\partial\phi_0}\right]V_{\text{eff}}(\phi_0) = 0. \tag{2.15}$$

Note that now we cannot use perturbation theory to determine $\beta(\lambda_0)$ and $\gamma(\lambda_0)$. As in the previous case, the effective potential displays a minimum at

$$\begin{gathered} 3\tilde{\lambda}v_0^2 = \frac{\Lambda^2}{\sqrt{e}}\exp\left[-\frac{16\pi^2}{9\tilde{\lambda}}\right], \\ V_{\text{eff}}(v_0) = -\frac{m_H^4}{128\pi^2}, \qquad m_H^2 = \omega^2(v_0). \end{gathered} \tag{2.16}$$

Using (2.15) at the minimum v_0, we get

$$\left[\Lambda\frac{\partial}{\partial\Lambda} + \beta(\lambda_0)\frac{\partial}{\partial\lambda_0}\right]m_H^2 = 0, \tag{2.17}$$

which in turn gives

$$\beta(\lambda_0) = -a_1\frac{9}{8\pi^2}\tilde{\lambda}^2. \tag{2.18}$$

This last equation implies that the theory is free asymptotically for $\Lambda \to \infty$ in agreement with triviality:

$$\tilde{\lambda} \sim \frac{16\pi^2}{9a_1} \frac{1}{\ln(\Lambda^2/m_H^2)}. \tag{2.19}$$

Inserting now (2.18) into (2.15), we obtain

$$\gamma(\lambda_0) = a_1^2 \frac{9}{16\pi^2} \tilde{\lambda}. \tag{2.20}$$

This last equation assures that $\tilde{\lambda}\phi_0^2$ is a renormalization group invariant. Rewriting the effective potential as

$$V_{\text{eff}}(\phi_0) = \frac{\left(3\tilde{\lambda}\phi_0^2\right)^2}{64\pi^2} \left[\ln\left(\frac{3\tilde{\lambda}\phi_0^2}{m_H^2}\right) - \frac{1}{2}\right], \tag{2.21}$$

we see that V_{eff} is manifestly renormalization group invariant.

Let us introduce the renormalized field η_R and condensate ϕ_R. Since the fluctuation η is a free field, we have $\eta_R = \eta$, namely

$$Z_\eta = 1. \tag{2.22}$$

On the other hand, for the scalar condensate, according to (2.20) we have

$$\phi_R = Z_\phi^{-1/2} \phi_0, \qquad Z_\phi \sim \lambda_0^{-1} \sim \ln\left(\frac{\Lambda}{m_H}\right). \tag{2.23}$$

As a consequence, we get that the physical mass m_H is *finitely* related to the renormalized vacuum expectation scalar field value v_R:

$$m_H = \xi v_R. \tag{2.24}$$

It should be clear that the physical mass m_H is an arbitrary parameter of the theory (dimensional transmutation). On the other hand, the parameter ξ being a pure number can be determined in the nonperturbative lattice approach.

3. The Higgs Boson Mass

The lattice approach to quantum field theories offers us the unique opportunity to study a quantum field theory by means of nonperturbative methods. Starting from the classical Lagrangian (2.1), one obtains the lattice theory defined by the Euclidean action:

$$S = \sum_x \left[\frac{1}{2} \sum_{\hat{\mu}} \left(\Phi(x+\hat{\mu}) - \Phi(x) \right)^2 + \frac{r_0}{2} \Phi^2(x) + \frac{\lambda_0}{4} \Phi^4(x) \right], \tag{3.1}$$

where x denotes a generic lattice site and, unless otherwise stated, lattice units are understood. It is customary to perform numerical simulations in the so-called Ising limit. The Ising limit corresponds to $\lambda_0 \to \infty$. In this limit, the one-component scalar field theory becomes governed by the lattice action

$$S_{\text{Ising}} = -\kappa \sum_x \sum_\mu \left[\phi(x+\hat{e}_\mu)\phi(x) + \phi(x-\hat{e}_\mu)\phi(x) \right] \tag{3.2}$$

with $\Phi(x) = \sqrt{2\kappa}\, \phi(x)$ and where $\phi(x)$ takes only the values +1 or −1.

It is known that there is a critical coupling [13]:

$$\kappa_c = 0.074834(15), \tag{3.3}$$

such that for $\kappa > \kappa_c$ the theory is in the broken phase, while, for $\kappa < \kappa_c$, it is in the symmetric phase. The continuum limit corresponds to $\kappa \to \kappa_c$ where $m_{\text{latt}} \equiv am_H \to 0$, a being the lattice spacing.

As discussed in Section 1, the triviality of the scalar theory means that the renormalized self-coupling vanishes as $1/\ln(\Lambda^2/m_H^2)$ when $\Lambda \to \infty$. As a consequence, in the continuum limit the theory admits a Gaussian fixed point.

On the lattice, the ultraviolet cutoff is $\Lambda = \pi/a$ so that we have

$$\lambda \sim \frac{1}{\ln(\Lambda/m_H)} \sim \frac{1}{\ln(\pi/am_H)} = \frac{1}{\ln(\pi/m_{\text{latt}})}. \tag{3.4}$$

The perturbative interpretation of triviality [4, 5] assumes that, in the continuum limit, there is an *infrared* Gaussian fixed point where the limit $m_{\text{latt}} \to 0$ corresponds to $m_H \to 0$. On the other hand, according to Section 2, in the triviality and spontaneous symmetry breaking scenario, the continuum dynamics is governed by an *ultraviolet* Gaussian fixed point where $m_{\text{latt}} \to 0$ corresponds to $a \to 0$. As we discuss below, these two different interpretations of triviality lead to different logarithmic correction to the Gaussian scaling laws which can be checked with numerical simulations on the lattice.

In [10], extensive numerical lattice simulations of the one-component scalar field theory in the Ising limit have been performed. In particular, using the Swendsen-Wang [14] and Wolff [15] cluster algorithms, the bare magnetization (vacuum expectation value):

$$v_{\text{latt}} = \langle |\phi| \rangle, \qquad \phi \equiv \frac{1}{L^4} \sum_x \phi(x), \tag{3.5}$$

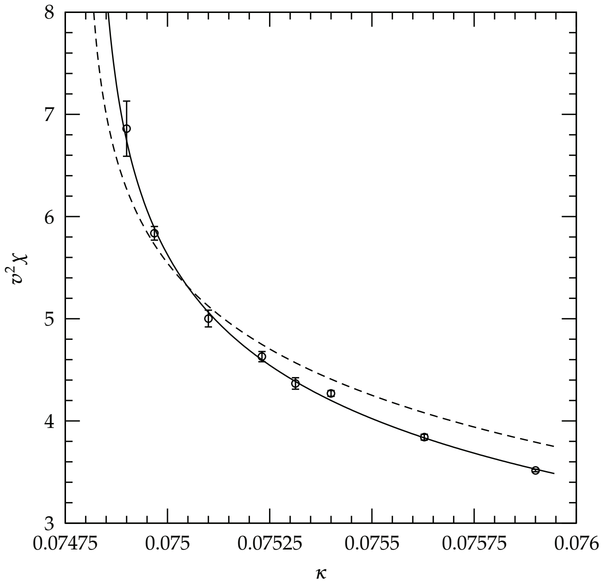

Figure 1: We show the lattice data for $v^2_{\mathrm{latt}}\chi_{\mathrm{latt}}$ together with the fit (3.9) (solid line) and the two-loop fit (3.11) (dashed line) where the fit parameters a_1 and a_2 are allowed to vary inside their theoretical uncertainties (3.12).

and the bare zero-momentum susceptibility:

$$\chi_{\mathrm{latt}} = L^4\left[\left\langle|\phi|^2\right\rangle - \langle|\phi|\rangle^2\right], \tag{3.6}$$

have been computed. According to the perturbative scheme of [4, 5], one expects

$$v^2_{\mathrm{latt}}\chi_{\mathrm{latt}} \sim |\ln(\kappa - \kappa_c)|, \quad \kappa \longrightarrow \kappa_c^+. \tag{3.7}$$

On the other hand, since, in the triviality and spontaneous symmetry breaking scenario, one expects that $Z_\varphi \sim \ln(\Lambda/m_H) \sim |\ln(\kappa - \kappa_c)|$, we have

$$v^2_{\mathrm{latt}}\chi_{\mathrm{latt}} \sim |\ln(\kappa - \kappa_c)|^2, \quad \kappa \longrightarrow \kappa_c^+. \tag{3.8}$$

The predictions in (3.8) can be directly compared with the lattice data reported in [10] and displayed in Figure 1. We fitted the data to the 2-parameter form:

$$v^2_{\mathrm{latt}}\chi_{\mathrm{latt}} = \alpha|\ln(\kappa - \kappa_c)|^2. \tag{3.9}$$

We obtain a rather good fit of the lattice data (full line in Figure 1) with

$$\alpha = 0.07560(49), \qquad \kappa_c = 0.074821(12), \qquad \chi^2_{\mathrm{dof}} \simeq 1.5. \tag{3.10}$$

Note that our precise determinations of the critical coupling κ_c in (3.10) are in good agreement with the value obtained in [13] (see (3.3)).

On the other hand, the prediction based on 2-loop renormalized perturbation theory is [5, 16] ($l = |\ln(\kappa - \kappa_c)|$):

$$\left[v_{\mathrm{latt}}^2 \chi_{\mathrm{latt}}\right]_{2\text{-loop}} = a_1 \left(l - \frac{25}{27} \ln l\right) + a_2 \tag{3.11}$$

together with the theoretical relations:

$$a_1 = 1.20(3), \qquad a_2 = -1.6(5). \tag{3.12}$$

We fitted the lattice data to (3.11) by allowing the fit parameters a_1 and a_2 to vary inside their theoretical uncertainties (3.12). The fit resulted in (dashed line in Figure 1)

$$a_1 = 1.17, \qquad a_2 = -2.10, \qquad \kappa_c = 0.074800(1), \qquad \chi^2_{\mathrm{dof}} \simeq 132. \tag{3.13}$$

It is evident from Figure 1 that the quality of the 2-loop fit is poor. However, these results have been criticized by the authors of [16] and have given rise to an intense debate in the recent literature [17–21].

Additional numerical evidences would come from the direct detection of the condensate rescaling $Z_\phi \sim |\ln(\kappa - \kappa_c)|$ on the lattice. To this end, we note that

$$Z_\phi \equiv 2\kappa m_{\mathrm{latt}}^2 \chi_{\mathrm{latt}}. \tag{3.14}$$

In Figure 2 we display the lattice data obtained in [10] for Z_ϕ, as defined in (3.14) versus m_{latt} reported in [5] at the various values of κ. For comparison, we also report the perturbative prediction of Z_η taken from [5]. We try to fit the lattice data with

$$Z_\phi = A \ln\left(\frac{\pi}{m_{\mathrm{latt}}}\right). \tag{3.15}$$

Indeed, we obtain a satisfying fit to the lattice data (solid line in Figure 2):

$$A = 0.498(5), \qquad \chi^2_{\mathrm{dof}} \simeq 4.1. \tag{3.16}$$

By adopting this alternative interpretation of triviality, there are important phenomenological implications. In fact, assuming to know the value of v_R, the ratio $\xi = m_H/v_R$ is now a cutoff-independent quantity. Indeed, the physical v_R has to be computed from the bare v_B through $Z = Z_\varphi$ rather than through the perturbative $Z = Z_\eta$. In this case, the perturbative relation [5]:

$$\frac{m_H}{v_R} \equiv \sqrt{\frac{\lambda_R}{3}}, \tag{3.17}$$

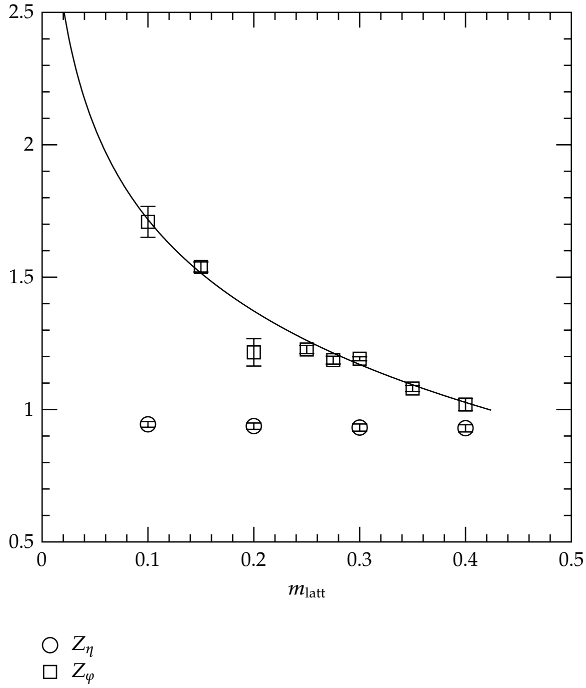

Figure 2: The lattice data for Z_ϕ, as defined in (3.14), and the perturbative prediction Z_η versus m_{latt}. The solid line is the fit to (3.15).

becomes

$$\frac{m_H}{v_R} = \sqrt{\frac{\lambda_R}{3}\frac{Z_\varphi}{Z_\eta}} \equiv \xi \tag{3.18}$$

obtained by replacing Z_η with Z_φ in [5] and correcting for the perturbative Z_η. Using the values of λ_R reported in [5] and our values of Z_φ, we display, in Figure 3, the values of m_H as defined through (3.18) versus m_{latt} for $v_R = 246\,\text{GeV}$. The error band corresponds to a one standard deviation error in the determination of m_H through a fit with a constant function. As one can see, the $Z_\varphi \sim \ln\Lambda$ trend observed in Figure 2 compensates the $1/\ln\Lambda$ from λ_R so that ξ turns out to be a cutoff-independent constant:

$$\xi = 3.065(80), \qquad \chi^2_{\text{dof}} \simeq 3.0, \tag{3.19}$$

which corresponds to

$$m_H = 754 \pm 20 \pm 20\,\text{GeV}, \tag{3.20}$$

where the last error is our estimate of systematic effects.

One could object that our lattice estimate of the Higgs mass (3.20) is not relevant for the physical Higgs boson. Indeed, the scalar theory relevant for the Standard Model is

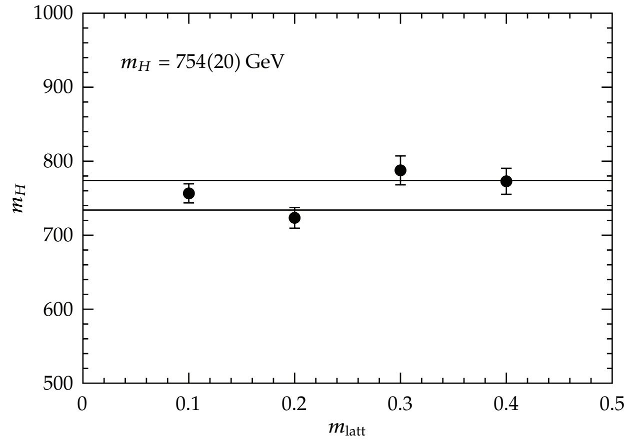

Figure 3: The values of m_H as defined through (3.18) versus $m_{\rm latt}$ assuming $v_R = 246\,\text{GeV}$. The error band corresponds to one standard deviation error in the determination of m_H through a fit with a constant function.

the O(4)-symmetric self-interacting theory. However, the Higgs mechanism eliminates three scalar fields leaving as physical Higgs field the radial excitation whose dynamics is described by the one-component self-interacting scalar field theory. Therefore, we are confident that our determination of the Higgs mass applies also to the Standard Model Higgs boson.

4. The Higgs Physics at LHC

Recently, both the ATLAS and CMS collaborations [22, 23] reported the experimental results for the search of the Higgs boson at the Large Hadron Collider running at $\sqrt{s} = 7\,\text{TeV}$, based on a total integrated luminosity between $1\,\text{fb}^{-1}$ and $2.3\,\text{fb}^{-1}$.

It is worthwhile to briefly discuss the main physical properties of our proposal for the trivial Higgs boson. For Higgs mass in the range $700 - 800\,\text{GeV}$, the main production mechanism at LHC is the gluon fusion $gg \to H$. The theoretical estimate of the production cross-section at LHC for centre of mass energy $\sqrt{s} = 7\,\text{TeV}$ is [24]

$$\sigma(gg \longrightarrow H) \simeq 0.06 - 0.14\,\text{pb}, \qquad 700\,\text{GeV} < m_H < 800\,\text{GeV}. \tag{4.1}$$

The gluon coupling to the Higgs boson in the Standard Model is mediated by triangular loops of top and bottom quarks. Since the Yukawa coupling of the Higgs particle to heavy quarks grows with quark mass, thus, balancing the decrease of the triangle amplitude, the effective gluon coupling approaches a nonzero value for large loop-quark masses. On the other hand, we argued that the Higgs condensate rescales with Z_ϕ. This means that if the fermions acquire a finite mass through the Yukawa couplings, then we are led to conclude that the coupling of the physical Higgs field to the fermions could be very different from the Standard Model Higgs boson. On the other hand, the coupling of the Higgs field to the gauge vector bosons is fixed by the gauge symmetries. So the coupling of our Higgs boson to the gauge vector bosons is the same as for the Standard Model Higgs boson. For large Higgs masses, the vector-boson fusion mechanism becomes competitive to gluon fusion Higgs production [24]:

$$\sigma\left(W^+W^- \longrightarrow H\right) \simeq 0.02 - 0.03\,\text{pb}, \qquad 700\,\text{GeV} < m_H < 800\,\text{GeV}. \tag{4.2}$$

The main difficulty in the experimental identification of a very heavy Standard Model Higgs ($m_H > 650\,\text{GeV}$) resides in the large width which makes impossible to observe a mass peak. However, in the triviality and spontaneous symmetry breaking scenario, the Higgs self-coupling vanishes so that the decay width is mainly given by the decays into pairs of massive gauge bosons. Since the Higgs is trivial, there are no loop corrections due to the Higgs self-coupling and we obtain, for the Higgs total width

$$\Gamma(H) \simeq \Gamma\left(H \longrightarrow W^+W^-\right) + \Gamma\left(H \longrightarrow Z^0Z^0\right), \tag{4.3}$$

where [1, 2]

$$\begin{aligned}
\Gamma\left(H \longrightarrow W^+W^-\right) &\simeq \frac{G_F m_H^3}{8\sqrt{2}\pi}\sqrt{1-4x_W}\left(1-4x_W+12x_W^2\right), \quad x_W = \frac{m_W^2}{m_H^2}, \\
\Gamma\left(H \longrightarrow Z^0Z^0\right) &\simeq \frac{G_F m_H^3}{16\sqrt{2}\pi}\sqrt{1-4x_Z}\left(1-4x_Z+12x_Z^2\right), \quad x_Z = \frac{m_Z^2}{m_H^2}.
\end{aligned} \tag{4.4}$$

Assuming $m_H \simeq 750\,\text{GeV}$, $m_W \simeq 80\,\text{GeV}$, and $m_Z \simeq 91\,\text{GeV}$, we obtain

$$\Gamma(H) \simeq 340\,\text{GeV}. \tag{4.5}$$

A thorough discussion of the experimental signatures of our trivial Higgs is presented in [25] where we compare our proposal with the recent data from ATLAS and CMS collaborations based on a total integrated luminosity between $1\,\text{fb}^{-1}$ and $2.3\,\text{fb}^{-1}$. In fact, we argue that the available experimental data seem to be consistent with our scenario.

5. Conclusions

The Standard Model requires the existence of a scalar Higgs boson to break electroweak symmetry and provide mass terms to gauge bosons and fermion fields. Usually the spontaneous symmetry breaking in the Standard Model is implemented within the perturbation theory which leads to predict that the Higgs boson mass squared is proportional to the self-coupling. However, there exist several results which point to vanishing scalar self-coupling. Therefore, within the perturbative approach, scalar field theories represent just an effective description valid only up to some cutoff scale, for without a cutoff, there would be no scalar self-interactions and without them no symmetry breaking. In other words, spontaneous symmetry breaking is incompatible with strictly local scalar fields in the perturbative approach.

In this paper, we have shown that local scalar fields are compatible with spontaneous symmetry breaking. In this case, the continuum dynamics is governed by an ultraviolet

Gaussian fixed point (triviality) and a nontrivial rescaling of the scalar condensate. We argued that nonperturbative lattice simulations are consistent with this scenario. Moreover, we find that the Higgs boson is rather heavy. Finally, the nontrivial rescaling of the Higgs condensate suggests that the whole issue of generation of fermion masses through the Yukawa couplings must be reconsidered.

References

[1] J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, *The Higgs Hunter's Guide*, Addison-Wesley, Redwood City, Calif, USA, 1990.

[2] A. Djouadi, "The anatomy of electroweak symmetry breaking. Tome I: the Higgs boson in the Standard Model," *Physics Reports*, vol. 457, no. 1–4, pp. 1–216, 2008.

[3] K. G. Wilson and J. Kogut, "The renormalization group and the ε expansion," *Physics Reports*, vol. 12, no. 2, pp. 75–199, 1974.

[4] M. Lüscher and P. Weisz, "Scaling laws and triviality bounds in the lattice $\varphi 4$ theory. (I). One-component model in the symmetric phase," *Nuclear Physics*, vol. 290, no. C, pp. 25–60, 1987.

[5] M. Lüscher and P. Weisz, "Scaling laws and triviality bounds in the lattice ϕ^4 theory. (II). One-component model in the phase with spontaneous symmetry breaking," *Nuclear Physics*, vol. 295, no. 1, pp. 65–92, 1988.

[6] R. Fernandez, J. Fröhlich, and A. D. Sokal, *Random Walks, Critical Phenomena, and Triviality in Quantum Field Theory*, Springer, Berlin, Germany, 1992.

[7] C. B. Lang, *Computer Stochastics in Scalar Quantum Field Theory*, arXiv:hep-lat/9312004v1, 1993.

[8] M. Consoli and P. M. Stevenson, "The non-trivial effective potential of the 'trivial' $\lambda\Phi 4$ theory: a lattice test," *Zeitschrift für Physik C*, vol. 63, no. 3, pp. 427–436, 1994.

[9] P. Cea, M. Consoli, and L. Cosmai, "Indications on the Higgs boson mass from lattice simulations," *Nuclear Physics B - Proceedings Supplements*, vol. 129-130, pp. 780–782, 2004.

[10] P. Cea, M. Consoli, and L. Cosmai, "Large rescaling of the Higgs condensate: theoretical motivations and lattice results," *Nuclear Physics B - Proceedings Supplements*, vol. 83-84, no. 1–3, pp. 658–660, 2000.

[11] S. Coleman and E. Weinberg, "Radiative corrections as the origin of spontaneous symmetry breaking," *Physical Review D*, vol. 7, no. 6, pp. 1888–1910, 1973.

[12] R. Jackiw, "Functional evaluation of the effective potential," *Physical Review D*, vol. 9, no. 6, pp. 1686–1701, 1974.

[13] D. S. Gaunt, M. F. Sykes, and S. McKenzie, "Susceptibility and fourth-field derivative of the spin-1/2 Ising model for T>Tc and d=4," *Journal of Physics A: General Physics*, vol. 12, no. 6, article 018, pp. 871–877, 1979.

[14] R. H. Swendsen and J.-S. Wang, "Nonuniversal critical dynamics in Monte Carlo simulations," *Physical Review Letters*, vol. 58, no. 2, pp. 86–88, 1987.

[15] U. Wolff, "Collective monte Carlo updating for spin systems," *Physical Review Letters*, vol. 62, no. 4, pp. 361–364, 1989.

[16] J. Balog, A. Duncan, R. Willey, F. Niedermayer, and P. Weisz, "The 4d one component lattice φ^4 model in the broken phase revisited," *Nuclear Physics B*, vol. 714, no. 3, pp. 256–268, 2005.

[17] P. Cea, M. Consoli, and L. Cosmai, *Large Logarithmic Rescaling of the Scalar Condensate: A Subtlety with Substantial Phenomenological Implications*, arXiv:hep-lat/0501013v1, 2005.

[18] P. M. Stevenson, "Comparison of perturbative RG theory with lattice data for the 4d Ising model," *Nuclear Physics B*, vol. 729, no. 3, pp. 542–557, 2005.

[19] J. Balog, F. Niedermayer, and P. Weisz, "Repairing Stevenson's step in the 4d Ising model," *Nuclear Physics B*, vol. 741, no. 3, pp. 390–403, 2006.

[20] P. Castorina, M. Consoli, and D. Zappalà, "An alternative heavy Higgs mass limit," *Journal of Physics G: Nuclear and Particle Physics*, vol. 35, no. 7, Article ID 072010, 2008.

[21] P. H. Lundow and K. Markstrom, "private communication," 2009.

[22] A. Nisati, "Higgs searches at ATLAS," in *Proceedings of the 25 International Symposium on Lepton Photon Interactions at High Energies*, Mumbai, India, August 2011.

[23] V. Sharma, "Higgs searches at CMS," in *Proceedings of the 25 International Symposium on Lepton Photon Interactions at High Energies*, Mumbai, India, August 2011.

[24] S. Dittmaier, *Handbook of LHC Higgs Cross Sections: 1. Inclusive Observables*, arXiv:1109.5922, 2011.

[25] P. Cea and L. Cosmai, *The Trivial Higgs at LHC*, arXiv:1101.0593v3, 2011.