The $H_T$ Higgs boson at the LHC Run 2

Paolo Cea

INFN - Sezione di Bari, Via Amendola 173 - 70126 Bari, Italy

Abstract

We further elaborate on the proposal that the Higgs boson should be a broad heavy resonance, referred to as true Higgs $H_T$, with mass around $750 \text{ GeV}$. We stress once again that within the Standard Model the true Higgs is the unique possibility to implement the spontaneous symmetry breaking of the local gauge symmetry by elementary, relativistic and strictly local scalar fields. We discuss the most relevant decay modes of the $H_T$ boson and estimate their partial decay widths and branching ratios. We discuss briefly the experimental signatures of the $H_T$ Higgs boson and compare with the recent available LHC data at $\sqrt{s} = 13 \text{ TeV}$. We find that the coupling of the $H_T$ Higgs boson to fermions is strongly suppressed. We also compare our theoretical expectations in the so-called golden channel to the data collected by the ATLAS and CMS Collaborations at $\sqrt{s} = 13 \text{ TeV}$ with an integrated luminosity of $36.1 \text{ fb}^{-1}$ and $35.9 \text{ fb}^{-1}$ respectively. We find that our theoretical expectations are in fair good agreement with the experimental observations. Combining the data from both the LHC Collaborations we obtain an evidence of the heavy Higgs boson in this channel with an estimated statistical significance of more than three standard deviations. Finally, we argue that, if the signal is real, by the end of the Run 2 both the LHC experiments will reach in the golden channel a statistical significance of about five standard deviations.

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1Electronic address: Paolo.Cea@ba.infn.it
1 Introduction

A fundamental feature of the Standard Model is the mechanism of spontaneous symmetry breaking, the so-called Brout-Englert-Higgs (BEH) mechanism [1, 2]. The first runs of proton-proton collisions at the CERN Large Hadron Collider (LHC) with center-of-mass energies $\sqrt{s} = 7, 8$ TeV (Run 1) has brought the confirmation of the existence of a boson, named H, which resembles the one which breaks the electroweak symmetry in the Standard Model of particle physics [5, 6]. The combined ATLAS and CMS experiments best estimate of the mass of the H boson was $m_H = 125.09 \pm 0.24$ GeV [7]. Moreover, the results from both LHC experiments, as summarized in Refs. [8, 9, 10], showed that all measurements of the properties of the new H resonance were consistent with those expected for the Standard Model Higgs boson. Actually, in the LHC Run 1 the strongest signal significance has been obtained from the decays of the H boson into two vector bosons, $H \rightarrow VV$ where $V = \gamma, W, Z$. In fact, in these channels the observed signal significance were above five standard deviations [8, 9, 10]. Nevertheless, soon after the evidence of the LHC resonance at 125 GeV, we already proposed [11] that the H resonance could be interpreted as a pseudoscalar meson. In particular, we showed [12] that our pseudoscalar meson could mimic the decays of the Standard Model Higgs boson in the vector boson decay channels, while the decays into fermions were strongly suppressed. Moreover, the main decay channels of this pseudoscalar meson were the hadronic decays that could mimic the decays of a putative Higgs boson in the hadronic channels. We feel that the only save way to distinguish experimentally our pseudoscalar meson from the Standard Model Higgs boson is to determine the $CP$ assignment of the H resonance. The spin and $CP$ properties of the $H$ boson can be determined by studying the tensor structure of its interactions with the electroweak gauge bosons. The Run 1 experimental analyses relied on discriminant observables chosen to be sensitive to the spin and parity of the signal. In this way it was possible to compare the Standard Model hypothesis $J^{PC} = 0^{++}$ to several alternative spin and parity models. It turned out that all tested alternative models were excluded with a statistical significance of about three standard deviations [13, 14]. In particular, spin-one and spin-two hypotheses were excluded at a 99 % CL or higher. Given the exclusion of the spin-one and spin-two scenarios, constraints were set on the anomalous couplings of the $H$ resonance to vector bosons by assuming a spin-zero state. Under the hypothesis that the new resonance is a spin-zero boson, the tensor structure of the interactions of the $H$ boson with two vector bosons were investigated and limits on CP-odd anomalous contributions were set. As a result, in the LHC Run 1 the pseudoscalar hypothesis was excluded at 99 % CL from CMS [14] and at 97.8 % CL from ATLAS [13]. However, this conclusion was not strengthened by the recent data from Run 2. As a matter of fact, the CMS Collaboration performed the study of the anomalous interactions of the $H$ resonance by using the full dataset recorded during the Run 2 corresponding to an integrated luminosity of 35.9 $fb^{-1}$ at $\sqrt{s} = 13$ TeV. Even tough the number of analyzed $H$ boson in Run 2 was about three times larger than in Run 1, the data were consistent with an almost equal mixture of scalar and pseudoscalar couplings [15]. Similar conclusions have been reached also by the ATLAS Collaboration [16] after collecting 36.1 $fb^{-1}$ in the Run 2. As a consequence, we may affirm that our pseudoscalar interpretation of the $H$ resonance cannot be yet completely excluded. Aside from these phenomenological considerations, we believe that there are different and compelling theoretical motivations to doubt on the identification of the $H$
resonance with the Standard Model Higgs boson. In fact, stemming from the known triviality problem, i.e. vanishing of the self-coupling, that affects self-interacting local scalar quantum fields in four space-time dimensions [17], it was evidenced that the Higgs boson condensation triggering the spontaneous breaking of the local gauge symmetries needs to be dealt with non-perturbatively. If this is the case, from one hand there is no stability problem for the condensate ground state, on the other hand the Higgs mass is finitely related to the vacuum expectation value of the quantum scalar field and it can be evaluated from first principles. Precise non-perturbative numerical simulations indicated that the true Higgs boson, henceforth denoted as $H_T$, is a rather heavy resonance with mass around $750 \text{ GeV}$ [18].

The aim of the present paper is to elaborate on the phenomenological consequences of the massive Higgs boson proposal. In particular, we will discuss the couplings of the $H_T$ Higgs boson to the massive vector bosons and to fermions, the expected production mechanism, and the main decay modes. We organize the paper as following. In Sect. 2, for sake of completeness, we briefly illustrate how spontaneous symmetry breaking arises in field theories involving scalar fields without quartic self-couplings. Section 3 is devoted to the discussion of the couplings of our massive Higgs boson proposal to Standard Model gauge fields. We determine the main decay channels of the $H_T$ Higgs boson and we critically examine the couplings to fermions. We also illustrate the main production mechanisms of the $H_T$ Higgs boson and estimate the production cross sections at the proton-proton collider at center-of-mass energy $\sqrt{s} = 13 \text{ TeV}$. In Sect. 4 we compare our proposal with available LHC Run 2 data from both ATLAS and CMS Collaborations. In Sect. 4.1 we try a quantitative comparison in the so-called golden channel of our theory with the recent data collected by the ATLAS and CMS Collaborations at $\sqrt{s} = 13 \text{ TeV}$ corresponding to an integrated luminosity of 36.1 fb$^{-1}$ and 35.9 fb$^{-1}$ respectively. Finally, our concluding remarks are relegated to Sect. 5.

2 Triviality and spontaneous symmetry breaking

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 $\lambda v^2$, where $\lambda$ is the renormalized scalar self-coupling and $v \sim 246 \text{ GeV}$ is the known weak scale. On the other hand, it is known that, within the non-perturbative description of spontaneous symmetry breaking in the Standard Model, self-interacting scalar fields are subject to the triviality problem [17], namely the renormalized self-coupling $\lambda \to 0$ when the ultraviolet cutoff is sent to infinity. Strictly speaking, there are no rigorous proof of triviality. Nevertheless, there exist several numerical studies which leave little doubt on the triviality conjecture. As a consequence, within the perturbative approach, the scalar sector of the Standard Model represents just an effective description valid only up to some cut-off scale. If the renormalized self-coupling of the scalar fields vanishes, then one faces with the problem of the spontaneous symmetry breaking mechanism and the related scalar Higgs boson. In fact, naively, one expects that the spontaneous symmetry breaking mechanism cannot be implemented without the scalar self-coupling $\lambda$. However, in Ref. [18], by means of nonperturbative numerical simulations of the $\lambda \Phi^4$ theory on the lattice, it was enlightened the scenario where the Higgs boson without self-interaction could coexist with spontaneous symmetry breaking. Moreover, due to the peculiar rescal-
ing of the Higgs condensate, the relation between the Higgs mass and $v$ is not the same as in perturbation theory. In fact, remarkably, it turned out that the Higgs mass were finitely related to $v$.

For reader convenience, in the present Section, following Ref. [18], we shall illustrate how spontaneous symmetry breaking could be compatible with triviality. To this end, we consider the simplest scalar field theory, namely a massless real scalar field $\Phi$ with quartic self-interaction $\lambda$ in four space-time dimensions:

$$L = \frac{1}{2}(\partial_\mu \Phi_0)^2 - \frac{1}{4}\lambda_0 \Phi_0^4,$$  \hspace{1cm} (2.1)

where $\lambda_0$ and $\Phi_0$ are the bare coupling and field respectively. As it is well known [19, 20], in the one-loop approximation the effective potential is given by:

$$V_{1-loop}^{eff}(\phi_0) = \frac{1}{4}\lambda_0 \phi_0^4 + \frac{1}{2} \int \frac{d^3k}{(2\pi)^3} \sqrt{\vec{k}^2 + 3\lambda_0 \phi_0^2}.$$  \hspace{1cm} (2.2)

This last equation shows that the one-loop effective potential is given the vacuum energy of the shifted field in the quadratic approximation. In fact, let us write:

$$\Phi_0 = \phi_0 + \eta$$  \hspace{1cm} (2.3)

where $\phi_0$ is the bare uniform scalar condensate, then in this approximation the Hamiltonian of the fluctuation $\eta$ over the background $\phi_0$ is:

$$\hat{H}_0 = \frac{1}{2}(\Pi_\eta)^2 + \frac{1}{2}(\nabla \eta)^2 + \frac{1}{2} (3\lambda_0 \phi_0^2) \eta^2 + \frac{1}{4}\lambda_0 \phi_0^4.$$  \hspace{1cm} (2.4)

After canonical quantization of the quadratic Hamiltonian $\hat{H}_0$ one readily finds that the energy density of the quantum vacuum in presence of the condensate $\phi_0$ is given by Eq. (2.2). It is, now, easy to see that the one-loop effective potential Eq. (2.2) displays a non-trivial minimum implying spontaneous symmetry breaking. However, the minimum of the effective potential lies outside the expected range of validity of the one-loop approximation and, therefore, it must be rejected as an artifact of the approximation [19, 20].

On the other hand, the triviality hypothesis implies that the fluctuation field $\eta$ is a free field with mass $\omega(\phi_0)$. As a consequence the exact effective potential is:

$$V_{eff}(\phi_0) = \frac{1}{4}\lambda_0 \phi_0^4 + \frac{1}{2} \int \frac{d^3k}{(2\pi)^3} \sqrt{\vec{k}^2 + \omega^2(\phi_0)} = \frac{1}{4}\lambda_0 \phi_0^4 + \frac{\omega^4(\phi_0)}{64\pi^2} \ln \left( \frac{\omega^2(\phi_0)}{\Lambda^2} \right),$$  \hspace{1cm} (2.5)

where $\Lambda$ is an ultraviolet cutoff. Moreover, the mechanism of spontaneous symmetry breaking implies that the mass of the fluctuation is related to the scalar condensate as:

$$\omega^2(\phi_0) = 3 \tilde{\lambda} \phi_0^2, \quad \tilde{\lambda} = a_1 \lambda_0, \quad (2.6)$$

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:

$$\left[ \Lambda \frac{\partial}{\partial \Lambda} + \beta(\lambda_0) \frac{\partial}{\partial \lambda_0} + \gamma(\lambda_0) \phi_0 \frac{\partial}{\partial \phi_0} \right] V_{eff}(\phi_0) = 0.$$  \hspace{1cm} (2.7)

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

$$3\tilde{\lambda}v_0^2 = \frac{\Lambda^2}{\sqrt{e}} \exp\left[-\frac{16\pi^2}{9\lambda}\right], \quad (2.8)$$

and

$$V_{eff}(v_0) = -\frac{m_{H_T}^4}{128\pi^2}, \quad m_{H_T}^2 = \omega^2(v_0). \quad (2.9)$$

Using Eq. (2.7) at the minimum $v_0$ we get:

$$\left[\Lambda \frac{\partial}{\partial \Lambda} + \beta(\lambda_0) \frac{\partial}{\partial \lambda_0}\right] m_{H_T}^2 = 0, \quad (2.10)$$

which in turns gives:

$$\beta(\lambda_0) = -a_1 \frac{9}{8\pi^2} \tilde{\lambda}^2. \quad (2.11)$$

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

$$\tilde{\lambda} \sim \frac{16\pi^2}{9a_1 \ln(\frac{\Lambda^2}{m_{H_T}})}.$$

(2.12)

Inserting now Eq. (2.11) into Eq. (2.7) we obtain:

$$\gamma(\lambda_0) = a_1^2 \frac{9}{16\pi^2} \tilde{\lambda}. \quad (2.13)$$

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

$$V_{eff}(\phi_0) = \frac{(3\tilde{\lambda} \phi_0^2)^2}{64\pi^2} \left[\ln\left(\frac{3\tilde{\lambda} \phi_0^2}{m_{H_T}^2}\right) - \frac{1}{2}\right], \quad (2.14)$$

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

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

$$Z_\eta = 1. \quad (2.15)$$

On the other hand, for the scalar condensate according to Eq. (2.13) we have:

$$\phi_R = Z_\phi^{-\frac{1}{2}} \phi_0, \quad Z_\phi \sim \lambda_0^{-1} \sim \ln(\frac{\Lambda}{m_H}). \quad (2.16)$$

As a consequence we get that the physical mass $m_{H_T}$ is finitely related to the renormalized vacuum expectation scalar field value $v$:

$$m_{H_T} = \xi v. \quad (2.17)$$

It should be clear that the physical mass $m_{H_T}$ is an arbitrary parameter of the theory (dimensional transmutation). On the other hand the parameter $\xi$ being a pure number can
be determined in the non-perturbative lattice approach. Remarkably, extensive numerical simulations showed that the physical Higgs boson mass $m_{H_T}$ is finitely related to the renormalized vacuum expectation value $v$. Moreover, the extrapolation to the continuum limit of the ratio $m_{H_T}/v$ leads to the intriguing relation [18]:

$$\xi \simeq \pi ,$$

(2.18)

pointing to a rather massive $H_T$ boson, namely $m_{H_T} \simeq 750$ GeV.

### 3 The $H_T$ Higgs boson

In the previous Section we presented simple arguments to illustrate how the spontaneous symmetry breaking mechanism can be implemented also for real scalar fields without self-interaction. One could object that our treatment could not be relevant for the physical Higgs boson, for the scalar theory relevant for the Standard Model is 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 theoretical arguments can be reliably extended to the Standard Model Higgs boson.

In order to determine the phenomenological signatures of the massive $H_T$ Higgs boson we need to take care of the couplings with the gauge and fermion fields of the Standard Model. Actually, the coupling of the Higgs field to the gauge vector bosons is fixed by the gauge symmetries. Therefore the coupling of the $H_T$ Higgs boson to the gauge vector bosons is the same as in perturbation theory notwithstanding the non-perturbative Higgs condensation driving the spontaneous breaking of the gauge symmetries. Given the rather large mass of the $H_T$ Higgs boson, the main decay modes are the decays into two massive vector bosons (see, e.g., Refs. [21, 22]):

$$\Gamma(\not H_T \rightarrow W^+ W^-) \simeq \frac{G_F m_{H_T}^3}{8\pi \sqrt{2}} \sqrt{1 - \frac{4m_{W}^2}{m_{H_T}^2}} \left(1 - 4 \frac{m_{W}^2}{m_{H_T}^2} + 12 \frac{m_{W}^4}{m_{H_T}^4}\right)$$

(3.1)

and

$$\Gamma(\not H_T \rightarrow Z^0 Z^0) \simeq \frac{G_F m_{H_T}^3}{16\pi \sqrt{2}} \sqrt{1 - \frac{4m_{Z}^2}{m_{H_T}^2}} \left(1 - 4 \frac{m_{Z}^2}{m_{H_T}^2} + 12 \frac{m_{Z}^4}{m_{H_T}^4}\right).$$

(3.2)

On the other hand, it is known that for heavy Higgs the radiative corrections to the decay widths can be safely neglected [23, 24, 25].

The couplings of the $H_T$ Higgs boson to the fermions are given by the Yukawa couplings $\lambda_f$. Unfortunately, there are not reliable lattice non-perturbative simulations on the continuum limit of the Yukawa couplings. If we follow the perturbative approximation, then the fermion Yukawa couplings turn out to be proportional to the fermion mass, $\lambda_f = \sqrt{2} m_f/v$. In that case, for heavy Higgs the only relevant fermion coupling is the top Yukawa coupling $\lambda_t$. On the other hand, we cannot exclude that the couplings of the physical Higgs field to the fermions could be very different from perturbation theory. Indeed, the non-trivial rescaling of the Higgs condensate suggests that, if the fermions acquire a finite mass through the Yukawa couplings, then the coupling of the physical

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2A preliminary account on this matter has been presented in the second part of Ref. [12].
Higgs field to the fermions could be strongly suppressed. In this case, the whole issue of generation of fermion masses through the Yukawa couplings should be reconsidered. Therefore, to parametrize our ignorance on the Yukawa couplings of the $H_T$ Higgs boson we introduce the parameter:

$$\kappa = \frac{\lambda^2}{2 m_t^2} v^2. \tag{3.3}$$

Obviously, in perturbation theory we have $\kappa = 1$. Nevertheless, we shall also consider the case $\kappa \simeq 0$ corresponding to strongly suppressed fermion Yukawa couplings.

The width for the decays of the $H_T$ boson into a $t\bar{t}$ pairs is easily found \cite{21, 22}:

$$\Gamma(H_T \to t\bar{t}) \simeq \kappa \frac{3 G_F m_{H_T} m_t^2}{4 \pi \sqrt{2}} \left(1 - 4 \frac{m_t^2}{m_{H_T}^2}\right)^{3/2}. \tag{3.4}$$

So that, to a good approximation, the Higgs total width is given by:

$$\Gamma_{H_T} \simeq \Gamma(H_T \to W^+ W^-) + \Gamma(H_T \to Z^0 Z^0) + \Gamma(H_T \to t\bar{t}). \tag{3.5}$$

Our previous equations show that, in the high mass region $m_{H_T} \gtrsim 400 \text{ GeV}$ the total width depends strongly on the Higgs mass. The main difficulty in the experimental identification of a very heavy Higgs resides in the large width which makes almost impossible to observe a mass peak. In fact, the expected mass spectrum of our heavy Higgs should be proportional to the Lorentzian distribution. For a resonance with mass $M$ and total width $\Gamma$ the Lorentzian distribution is given by:

$$L(E) \sim \frac{\Gamma}{(E - M^2)^2 + \Gamma^2/4}. \tag{3.6}$$

Note that Eq. (3.6) is the simplest distribution consistent with the Heisenberg uncertainty principle and the finite lifetime $\tau \simeq 1/\Gamma$. We obtain, therefore, the following Lorentzian distribution:

$$L_{H_T}(E) \sim \frac{1}{1.0325 \pi} \frac{\Gamma_{H_T}(E)}{(E - 750 \text{ GeV})^2 + \left(\frac{\Gamma_{H_T}(E)}{2}\right)^2}, \tag{3.7}$$

where $\Gamma_{H_T}(E)$ is given by Eq. (3.3), and the normalization is such that:

$$\int_0^\infty L_{H_T}(E) \, dE = 1. \tag{3.8}$$

Note that, in the limit $\Gamma_{H_T} \to 0$, $L_{H_T}(E)$ reduces to $\delta(E - 750 \text{ GeV})$.

To evaluate the Higgs event production at LHC we need the inclusive Higgs production cross section. As in perturbation theory, for large Higgs masses the main production processes are by vector-boson fusion and gluon-gluon fusion. In fact, the $H_T$ Higgs production cross section by vector-boson fusion is the same as in the perturbative Standard Model calculations. Moreover, for Higgs mass in the range $700 - 800 \text{ GeV}$ the main production mechanism at LHC is expected to be by the gluon fusion mechanism. The gluon coupling to the Higgs boson in the Standard Model is mediated by triangular loops of top and bottom quarks. Since in perturbation theory the Yukawa couplings of the Higgs particle to heavy quarks grows with quark mass, thus balancing the decrease of the triangle
Figure 1: The inclusive Higgs production gluon-gluon fusion (left panel) and vector-boson fusion (right panel) cross sections at $\sqrt{s} = 13$ TeV. The data have been taken from Ref. [26]. The dashed lines are the fits to the data with our parametrization Eqs. (3.10) and (3.12).

amplitude, the effective gluon coupling approaches a non-zero value for large loop-quark masses. This means that for heavy Higgs the gluon fusion inclusive cross section is almost completely determined by the top quarks, even though it is interesting to stress that for large Higgs masses the vector-boson fusion mechanism becomes competitive to the gluon fusion Higgs production. According to our approximations the total inclusive cross section for the production of the $H_T$ Higgs boson can be written as:

$$\sigma(pp \rightarrow H_T + X) \simeq \sigma_{VV}(pp \rightarrow H_T + X) + \kappa \sigma_{gg}(pp \rightarrow H_T + X), \quad (3.9)$$

where $\sigma_{VV}$ and $\sigma_{gg}$ are the vector-boson fusion and gluon-gluon fusion inclusive cross sections respectively.

In Ref. [26] it is presented the calculations of the cross sections computed at next-to-next-to-leading and next-to-leading order for heavy Higgs boson with Standard Model-like coupling at $\sqrt{s} = 13$ TeV. In Fig. 1 we show the dependence on the Higgs mass of the perturbative gluon-gluon (left panel) and vector-boson fusion (right panel) cross sections at $\sqrt{s} = 13$ TeV as reported in Ref. [26]. As concern the gluon-gluon fusion cross section we found that this cross section can be parametrized as:

$$\sigma_{gg}(pp \rightarrow H_T + X) \simeq \begin{cases} \left( \frac{a_1}{M_{H_T}} + a_2 M_{H_T}^3 \right) \exp(-a_3 M_{H_T}) & M_{H_T} \leq 300 \text{ GeV} \\ a_4 \exp \left[ -a_5 (M_{H_T} - 400 \text{ GeV}) \right] & 300 \text{ GeV} \leq M_{H_T} \leq 400 \text{ GeV} \\ 400 \text{ GeV} \leq M_{H_T} \end{cases}, \quad (3.10)$$

where $M_{H_T}$ is expressed in GeV. In fact we fitted Eq. (3.10) to the data (see the dashed line in Fig. 1 left panel) and obtained:

$$a_1 \simeq 1.24 \times 10^4 \text{ pb GeV}, \quad a_2 \simeq 1.49 \times 10^{-6} \text{ pb GeV}^{-3},$$

$$a_3 \simeq 7.06 \times 10^{-3} \text{ GeV}^{-1}, \quad a_4 \simeq 9.80 \text{ pb},$$

$$a_5 \simeq 7.63 \times 10^{-3} \text{ GeV}^{-1}. \quad (3.11)$$

Likewise, we parametrized the dependence of the vector-boson fusion cross section as:

$$\sigma_{VV}(pp \rightarrow H_T + X) \simeq \left( b_1 + \frac{b_2}{M_{H_T}} + \frac{b_3}{M_{H_T}^2} \right) \exp(-b_4 M_{H_T}), \quad (3.12)$$
and obtained for the best fit (see the dashed line in Fig. 1 right panel):

\[ b_1 \simeq -2.69 \times 10^{-6} \text{ pb} , \quad b_2 \simeq 8.08 \times 10^{2} \text{ pb GeV} , \]
\[ b_3 \simeq -1.98 \times 10^{4} \text{ pb GeV}^2 , \quad b_4 \simeq 2.26 \times 10^{-3} \text{ GeV}^{-1} . \]

(3.13)

4 \quad H_T \text{ Decay Channels}

We have seen that the main decays of the heavy Higgs boson are the decays into two massive vector boson and tt pairs, if \( \kappa = 1 \). Since the search for the decays of the Higgs boson into pairs of top quarks is very complicated because of the large QCD background, we focus on the decays into two massive vector bosons. It is worthwhile to stress that the Higgs decay to \( \gamma \gamma \) do not proceed directly, but, to a fair good approximation, via longitudinal W virtual states. Therefore the relevant branching ratio turns out to be strongly suppressed, i.e. \( Br(H_T \to \gamma \gamma) \sim 10^{-6} \) \cite{27,28}.

To compare the invariant mass spectrum of our Higgs with the experimental data, we observe that:

\[ N_{H_T}(E_1, E_2) \simeq \mathcal{L} \int_{E_1}^{E_2} Br(E) \varepsilon(E) \sigma(pp \to H_T + X) L_{H_T}(E) dE , \]

(4.1)

where \( N_{H_T} \) is the number of Higgs events in the energy interval \( E_1, E_2 \), corresponding to an integrated luminosity \( \mathcal{L} \), in the given channel with branching ratio \( Br(E) \). The parameter \( \varepsilon(E) \) accounts for the efficiency of trigger, acceptance of the detectors, the kinematic selections, and so on. Thus, in general \( \varepsilon(E) \) depends on the energy, the selected channel and the detector. For illustrative purposes, we consider the decay channels \( H \to WW \to \text{fermions} \) and \( H \to ZZ \to \text{fermions} \). Thus, for the branching ratios \( Br(E) \) we can write:

\[ Br(H_T \to WW \to \text{fermions}) \simeq Br(H_T \to WW) \times Br(WW \to \text{fermions}) \]
\[ Br(H_T \to ZZ \to \text{fermions}) \simeq Br(H_T \to ZZ) \times Br(ZZ \to \text{fermions}) , \]

(4.2)

where

\[ Br(H_T \to WW) = \frac{\Gamma(H_T \to WW)}{\Gamma_{H_T}} , \quad Br(H_T \to ZZ) = \frac{\Gamma(H_T \to ZZ)}{\Gamma_{H_T}} , \]

(4.3)

while the branching ratios for the decays of W and Z bosons into fermions are given by the Standard Model values \cite{29}.

In Fig. 2 we show the observed limits for a new neutral spin-zero resonance decaying to two Z bosons. The limits on the inclusive production cross section times the relevant branching ratio (the dashed lines in Fig. 2) have been obtained by the ATLAS Collaboration from the analysis of the data collected during the LHC Run 2 with an integrated luminosity of 36.1 fb\(^{-1}\) \cite{30} by using the decay channels where the pair of Z bosons decay leptonically to \( \ell^+ \ell^- \ell^+ \ell^- \) or \( \ell^+ \ell^- \nu \bar{\nu} \) final states, \( \ell \) being either an electron or a muon. The displayed limits correspond to gluon-gluon fusion (left panel) and vector-boson fusion (right panel) production mechanisms by assuming a narrow width Higgs-like scalar resonance. This allows us to directly compare the data to our theoretical estimates of the inclusive production cross section, Eqs. (3.10) and (3.12) times the branching ratios as given by Eq. (4.3). In fact, in Fig. 2 the continuum lines are our theoretical cross-section
Figure 2: (color online) Observed limits on the cross-section times branching ratio to ZZ final states for a narrow-width heavy scalar resonance as a function of its mass. The limits correspond to the decays of the two Z vector bosons into $\ell^+\ell^-\ell^+\ell^-$ or $\ell^+\ell^-\nu\bar{\nu}$, where $\ell$ is an electron or muon. The limits correspond to gluon-gluon fusion (left panel) and vector-boson fusion (right panel) production mechanisms. The data (dashed green lines) have been obtained from Fig. 6 in Ref. [30]. The continuum black lines are our theoretical estimates of the cross-section times branching ratio for gluon-gluon fusion processes, Eqs. (3.10) and (3.11), and vector-boson fusion processes, Eq. (3.12) and (3.13).
Figure 3: (color online) Comparison to the LHC data of the distribution of the invariant mass $m_{ZZ}$ for the process $H_T \rightarrow ZZ \rightarrow \ell\ell\ell\ell$ ($\ell = e, \mu$) in the high-mass region $m_{ZZ} \gtrsim 600 \text{GeV}$. The CMS data (left panel) have been obtained from Fig. 3, left panel, in Ref. [31] and correspond to an integrated luminosity of $\mathcal{L} = 35.9 \text{fb}^{-1}$. The ATLAS data (right panel), corresponding to an integrated luminosity of $\mathcal{L} = 36.1 \text{fb}^{-1}$, have been obtained from Fig. 4, left panel, in Ref. [30]. The dashed (green) lines are our estimate of the background, the continuum (red) lines are the signal histogram assuming $\varepsilon(E) \simeq 0.80$ and $\kappa \simeq 0$.

4.1 The golden channel

In the present Section we attempt a direct comparison of our theoretical expectations with the available experimental data from LHC Run 2 in the so-called golden channel corresponding to the decays $H_T \rightarrow ZZ \rightarrow \ell\ell\ell\ell$. Indeed, the four-lepton channel, albeit rare, has the clearest and cleanest signature of all the possible Higgs boson decay modes due to the small background contamination. In Fig. 3 we show the invariant mass distribution for the golden channel obtained from the CMS experiment with an integrated luminosity of $35.9 \text{fb}^{-1}$ [31] (left panel) and the ATLAS experiment with an integrated luminosity of $36.1 \text{fb}^{-1}$ [30] (right panel). From our estimate of the background (dashed lines in Fig. 3) we see that, indeed, in the high invariant mass region $m_{ZZ} \gtrsim 650 \text{GeV}$, the background is strongly suppressed. To compare with our theoretical expectations, we displayed in Fig. 3 the distribution of the invariant mass for the $H_T$ Higgs boson candidates corresponding to the golden channel. The distributions have been obtained using Eq. (4.1) with $\kappa = 0$ and $\varepsilon(E) \simeq 0.80$ to take care of the fact that the detectors do not cover the full phase space. We also assumed a slightly smaller value for the heavy Higgs boson central mass, namely $m_{H_T} \simeq 730 \text{GeV}$ that, however, is within the statistical uncertainties of the lattice determination [18].

Due to the rather low integrated luminosity, we see that the signal manifest itself as a broad plateau in the invariant mass interval $650 \text{GeV} \lesssim m_{ZZ} \lesssim 1000 \text{GeV}$ over a smoothly decreasing background. Actually, in this region it seems that there is a moderate excess of signal over the expected background that seems to compare quite well with our theoretical prediction. To be qualitative, we may estimate the total number of events in the invariant mass interval $650 \text{GeV} \lesssim m_{ZZ} \lesssim 1000 \text{GeV}$ and compare with our
Figure 4: (color online) Comparison to the LHC data of the distribution of the invariant mass $m_{ZZ}$ in the high-mass region $m_{ZZ} \gtrsim 600\,\text{GeV}$ for the process $H_T \to ZZ \to \ell\ell\ell\ell$ ($\ell = e, \mu$). The signal distribution has been obtained from the combination of the ATLAS and CMS event distributions by subtracting the relevant background. The continuum (red) line is the expected signal histogram assuming $\varepsilon(E) \simeq 0.80$ and $\kappa \simeq 0$.

Theoretical expectations. We find:

\begin{align}
N_{\text{obs}} &= 8.0^{+5.00}_{-2.83} \, , \quad N_{\text{back}} = 1.73 \, , \quad N_{\text{sign}}^{\text{obs}} = 6.27^{+5.00}_{-2.83} \, , \quad N_{\text{sign}}^{\text{th}} = 5.93 \quad \text{CMS} \\
N_{\text{obs}} &= 23.0^{+8.14}_{-4.70} \, , \quad N_{\text{back}} = 11.0 \, , \quad N_{\text{sign}}^{\text{obs}} = 12.00^{+8.14}_{-4.70} \, , \quad N_{\text{sign}}^{\text{th}} = 5.96 \quad \text{ATLAS}
\end{align}

where, to be conservative, the quoted errors have been obtained by adding in quadrature the experimental errors. Remarkably, both ATLAS and CMS distributions display an excess over the expected background with a statistical significance of more than two standard deviations. Moreover, the excesses interpreted as a signal seem to be in fair agreement with our theoretical expectations. To better appreciate this point, in Fig. 4 we show the signal-distribution of the invariant mass $m_{ZZ}$ in the high-mass region $m_{ZZ} \gtrsim 600\,\text{GeV}$. The signal has been obtained by combining the ATLAS and CMS binned events and subtracting the relevant background. It is remarkable that the signal distribution displays a broad peak structure around $m_{ZZ} \sim 700\,\text{GeV}$ with a statistical significance well above three standard deviations. Moreover, we see that our theoretical signal distribution (continuum line in Fig. 4) is in reasonable agreements with the experimental data. Therefore, we may conclude that our proposal for the heavy $H_T$ Higgs boson is finding in the golden channel the first confirmation, even though we cannot yet completely exclude the compatibility of the data with the background-only hypothesis.

5 Conclusion

It is widely believed that the new LHC resonance at $125\,\text{GeV}$ is the Standard Model Higgs boson. However, stemming from the known triviality problem, i.e. vanishing self-coupling, that affects self-interacting scalar quantum fields in four space-time dimensions,
we evidenced that the Higgs boson condensation triggering the spontaneous breaking of
the local gauge symmetries needs to be dealt with non perturbatively. It is worthwhile to
notice that in this case, from one hand there is no stability problem for the condensate
ground state, on the other hand the Higgs mass is finitely related to the vacuum expec-
tation value of the quantum scalar field and, in principle, it can be evaluated from first
principles. In fact, precise non-perturbative numerical simulations \[18\] gave for the $H_T$
Higgs boson mass $m_{H_T} = 754 \pm 20 GeV$ leading to a rather heavy Higgs boson. In this
paper we elaborated some phenomenological aspects of the heavy Higgs boson scenario.
We have critically discussed the couplings of the $H_T$ Higgs boson to the massive vector
bosons and to fermions. We have also estimated the expected production mechanism and
the main decay modes. Comparing with the available LHC Run 2 data we concluded that
the coupling of the $H_T$ Higgs boson to fermions were strongly suppressed. Finally, we
compared our proposal with the recent results in the golden channel from both ATLAS
and CMS Collaborations. We found that the available experimental observations were
consistent with our scenario. We are confident that forthcoming data from LHC Run 2
will add further support to the heavy Higgs proposal. Indeed, assuming a real signal, by
the end of the LHC Run 2 it is expected that both ATLAS and CMS experiments will
reach in the golden channel a statistical significance of about five standard deviations.

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