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Enhanced diphoton Higgs decay rate and isospin symmetric Higgs boson

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The ATLAS and CMS experiments have recently discovered a new 125 GeV boson. We show that the properties of this particle, including the enhancement of its diphoton decay rate, can be explained in a model with an isospin symmetric Higgs boson. The predictions of the model relevant for future experiments are also discussed.

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I. INTRODUCTION

Recently, the ATLAS [1] and CMS [2] experiments at the Large Hadron Collider (LHC) reported that a new boson h, compatible to the Standard Model (SM) Higgs boson H, was discovered in the mass range 125–126 GeV. On the other hand, the ATLAS and CMS data might already suggest existence of a new physics beyond the SM: While the decay channels of $h \to ZZ^*$ and $h \to WW^*$ are fairly consistent with the SM, the diphoton branching ratio $\mathrm{Br}(h \to \gamma\gamma)$ is about 1.6 times larger than the SM value.[1] This deviation from the SM has been discussed by many authors [5].

In this paper, we will show that the ATLAS and CMS data for the enhanced diphoton branching ratio can be explained in the class of models with isospin symmetric (IS) electroweak Higgs boson suggested by the authors in Refs. [6,7]. It is noticeable that as will be shown below, these models also make several predictions, which can be checked at the LHC in the near future.

II. IS HIGGS MODELS

There is a large hierarchy between quark masses from different families [8]. Besides, the isospin violation in different families is also hierarchical. It is very strong in the third family, strong (although essentially weaker) in the second family, and mild in the first one: $\frac{m_t}{m_b} \simeq 41.5$, $\frac{m_c}{m_s} \simeq 13.4$, and $\frac{m_u}{m_d} \simeq 0.38$–$0.58$ [8]. This is a big mystery: in the framework of the SM, it is required to introduce hierarchical Yukawa couplings by hand, e.g., $\frac{y_t^{\mathrm{SM}}}{y_b^{\mathrm{SM}}} \simeq 41.5$, and $\frac{y_c^{\mathrm{SM}}}{y_s^{\mathrm{SM}}} \simeq 13.4$.

A class of models (the IS Higgs models) describing the hierarchies in the quark mass spectrum was previously studied in Refs. [6,7]. One of our main motivations for introducing such models was to find a dynamical mechanism that could shed light on the experimental fact that the isospin violation in the quark mass spectrum is essentially stronger in heavier families.

The main characteristics of these models are the following: (a) it is assumed that the dynamics primarily responsible for electroweak symmetry breaking (EWSB) leads to the mass spectrum of quarks with no (or weak) isospin violation. Moreover, it is assumed that the values of these masses are of the order of the observed masses of the down-type quarks. (b) the second (central) assumption is introducing the horizontal interactions for the quarks in the three families. As a first step, a *subcritical* (although near-critical, i.e., strong) diagonal horizontal interactions for the top quark is utilized which lead to the observed ratio $\frac{m_t}{m_b} \simeq 41.5$. The second step is introducing *equal* strength (i.e., isospin symmetric) horizontal flavor-changing-neutral (FCN) interactions between the t and c quarks and the b and s ones.

All together, these interactions naturally provide the observed ratio $m_c/m_s \simeq 13.4$ in the second family [6]. As emphasized in Ref. [6], the choice of the IS masses being close to the values of the observed masses of the down-type quarks is crucial in this scenario. As to the mild isospin violation in the first family, it was studied together with the effects of the family mixing, reflected in the Cabibbo-Kobayashi-Maskawa (CKM) matrix [6] (see also Sec. V below).

In this scenario, besides the EWSB interactions, the dominant dynamics responsible for the form of the mass spectrum of quarks is connected with the diagonal horizontal interactions for the third family and the horizontal, isospin symmetric, FCN interactions between the second and third ones. One of the signatures of this scenario is the appearance of a composite top-Higgs doublet Φ_{h_t} (resonance) composed of the quarks and antiquarks of the third family [6,7].[2]

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[1] To the contrary, Plehn and Rauch [3] have recently argued that none of the measured couplings deviates from its SM values significantly. Also, the QCD uncertainties are discussed in Ref. [4]. Thus, the observed deviations are not yet definitive.

[2] Such composites in the near-critical regime in a symmetric phase of models with dynamical chiral symmetry breaking were studied by several authors [9].

Thus, the main source of the isospin violation in this approach is only the strong top-quark interactions. On the other hand, because these interactions are subcritical, the top quark plays a minor role in EWSB. The latter distinguishes this scenario from the top quark condensate model [10–15]. Note that unlike the topcolor assisted technicolor model (TC2) [16], this class of models utilizes subcritical dynamics for the top quark, so that without strong fine-tuning, the bosons from the top-Higgs doublet Φ_{h_t} are heavy, say, of order 1 TeV, in general (compare with Ref. [7]).

Although the concrete model in Refs. [6,7] utilized the fourth family of fermions [17,18] for generating EWSB, this choice is not crucial, as the authors emphasized in Ref. [6]. In particular, the fourth family can be replaced by just a IS Higgs boson doublet Φ_h, without specifying its composite origin (if any). In this paper, we will consider just such a version in which the neutral scalar from the Φ_h doublet will be identified with the 125 GeV h boson. Here we emphasize that while the neutral top-Higgs boson h_t has a large top-Yukawa coupling, the IS Higgs boson h does not, $y_t \simeq y_b \sim 10^{-2}$. On the other hand, the hWW^* and hZZ^* coupling constants are close to those in the SM (see below). Also, the mixing between Φ_h and much heavier Φ_{h_t} should be small (compare with Ref. [7]). Let us now describe the decay processes of the IS Higgs h.

III. DECAY MODES $h \to \gamma\gamma$, $h \to Z\gamma$, $h \to WW^*$, AND $h \to ZZ^*$

Let us consider the diphoton branching ratio in the IS Higgs model. It is well known that the W-loop contribution to $H \to \gamma\gamma$ is dominant in the SM, while the top-loop effect is destructive against the W-loop. More concretely, the diphoton partial width in the SM reads [19]

$$\Gamma(H \to \gamma\gamma) = \frac{\sqrt{2} G_F \alpha^2 m_H^3}{256\pi^3} \left| A_1(\tau_W) + N_c Q_t^2 A_{\frac{1}{2}}(\tau_t) \right|^2, \quad \tau_W \equiv \frac{m_H^2}{4m_W^2}, \qquad \tau_t \equiv \frac{m_H^2}{4m_t^2}, \tag{1}$$

where G_F denotes the Fermi constant, $N_c = 3$ represents the number of colors, and $Q_t = +2/3$ is the electric charge of the top quark. The loop functions A_1 and $A_{\frac{1}{2}}$ for W and t, respectively, are given by

$$A_1(\tau) \equiv -\frac{1}{\tau^2}[2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)], \tag{2}$$

and

$$A_{\frac{1}{2}}(\tau) \equiv \frac{2}{\tau^2}[\tau + (\tau - 1)f(\tau)], \tag{3}$$

with $f(\tau) \equiv \arcsin^2\sqrt{\tau}$ for $\tau \le 1$. Then, the numerical values of the W- and t-loop functions read

$$A_1(\tau_W) = -8.32, \qquad A_{\frac{1}{2}}(\tau_t) = 1.38, \tag{4}$$

for $m_W = 80.385$ GeV [8], $m_t = 173.5$ GeV [8], and $m_H = 125$ GeV.

On the other hand, in the IS Higgs model, the Yukawa coupling between the top and the IS Higgs h is as small as the bottom Yukawa coupling, so that the top-loop contribution is strongly suppressed. The partial decay width of $h \to \gamma\gamma$ is thus enhanced without changing essentially $h \to ZZ^*$ and $h \to WW^*$. A rough estimate taking the isospin symmetric top and bottom Yukawa couplings $y_t \simeq y_b \approx 10^{-2}$ is as follows:

$$\frac{\Gamma^{\rm IS}(h \to \gamma\gamma)}{\Gamma^{\rm SM}(H \to \gamma\gamma)} \simeq 1.56, \qquad \frac{\Gamma^{\rm IS}(h \to WW^*)}{\Gamma^{\rm SM}(H \to WW^*)} = \frac{\Gamma^{\rm IS}(h \to ZZ^*)}{\Gamma^{\rm SM}(H \to ZZ^*)} = \left(\frac{v_h}{v}\right)^2 \simeq 0.96. \tag{5}$$

Here using the Pagels-Stokar formula [20], we estimated the vacuum expectation value of the top-Higgs h_t as $v_t = 50$ GeV, and the vacuum expectation value v_h of the IS Higgs h is given by the relation $v^2 = v_h^2 + v_t^2$ with $v = 246$ GeV. Note that the values of the ratios in Eq. (5) are not very sensitive to the value of v_t, e.g., for $v_t = 40$–100 GeV, the suppression factor in the pair decay modes to WW^* and ZZ^* is $0.97 - 0.84$ and the enhancement factor in the diphoton channel is $1.58 - 1.37$. For the decay mode of $h \to Z\gamma$, this model yields

$$\frac{\Gamma^{\rm IS}(h \to Z\gamma)}{\Gamma^{\rm SM}(H \to Z\gamma)} \simeq 1.07 \tag{6}$$

(the data concerning this decay channel has not yet been reported [1,2]). Note that the total decay width is almost unchanged, so that Eqs. (5) and (6) indicate the suppression/enhancement factors of the corresponding branching ratios.

The values in Eq. (5) agree well with the data in the ATLAS and CMS experiments. However, obviously, the main production mechanism of the Higgs boson, the gluon fusion process $gg \to h$, is now in trouble. The presence of new chargeless colored particles, which considered by several authors [21] can help to resolve this problem. We pursue this possibility below.

IV. MODEL WITH COLORED SCALAR

We utilize an effective theory near the EWSB scale. The model contains: (1) the IS Higgs doublet Φ_h, which is mainly responsible for the EWSB and couples to the top and bottom in the isospin symmetric way, (2) the top-Higgs doublet Φ_{h_t}, which is required to obtain the correct top mass, and (3) the colored scalar and/or fermions which are required to enhance $gg \to h$.

The items (1) and (2) above are essentially described in Refs. [6,7]. The only difference is that the two composite Higgs doublets composed of the fourth family quarks should now be replaced by the IS Higgs. We will discuss this point later. Note that in this case the Lagrangian density $\mathcal{L}$ in the effective theory contains the IS Higgs quartic coupling λ_h, $\mathcal{L} \supset -\lambda_h|\Phi_h^\dagger \Phi_h|^2$, and the mass $m_h = 125$ GeV corresponds to a small λ_h via the relation $m_h^2 \simeq 2\lambda_h v_h^2$ like in the SM, because the mixing between Φ_h and the much heavier Φ_{h_t} is tiny in the present model (compare with Refs. [6,7]). However, unlike the case of the SM [22], this does not imply that the theory keeps the perturbative nature up to some extremely high energy scale, as we will see below.

As to a concrete realization of item (3), we may introduce a real scalar field S in the adjoint representation of the color $SU(3)_c$ and utilize the Higgs-portal model [21], just as a benchmark case,

$$\mathcal{L} \supset \mathcal{L}_S = \frac{1}{2}(D_\mu S)^2 - \frac{1}{2}m_{0,S}^2 S^2 - \frac{\lambda_S}{4}S^4 - \frac{\lambda_{hS}}{2}S^2\Phi_h^\dagger\Phi_h,$$
$$\Phi_h = \begin{pmatrix} \omega^+ \\ \frac{1}{\sqrt{2}}(v_h + h + iz_0) \end{pmatrix}, \tag{7}$$

where $\omega^\pm$ and z_0 are the components eaten by $W^\pm$ and Z. The scalar field S is chosen to be assigned to the $(\mathbf{8},\mathbf{1})_0$ representation of the $SU(3)_c \times SU(2)_W \times U(1)_Y$. Other representations, for example, a color triplet, are also possible. Note that we do not incorporate a Higgs-portal term between S and Φ_{h_t} and other possible cubic and quartic terms into Eq. (7), because they do not play any important role in the following analysis.

The mass-squared term for the scalar S is given by

$$M_S^2 = m_{0,S}^2 + \frac{\lambda_{hS}}{2}v_h^2, \tag{8}$$

and should be positive in order to avoid the color symmetry breaking. Typically, $M_S \sim 200$ GeV is allowed in the current data [21]. We will take a positive value for λ_{hS} and a classically (quasi-)scale invariant model with $m_{0,S}^2 \approx 0$, which is favorable to reproduce the SM like gluon fusion production.

Let us consider the contribution of the color octet S to the gluon fusion process $gg \to h$ in the leading order,

$$\frac{\sigma(gg\to h)}{\sigma^{\rm SM}(gg\to H)} \sim \frac{\Gamma(h\to gg)}{\Gamma^{\rm SM}(H\to gg)} = \left|\frac{C_A\lambda_{hS}\frac{vv_h}{2M_S^2}A_0(\tau_S)}{A_{\frac{1}{2}}(\tau_t)}\right|^2, \tag{9}$$

with $C_A = 3$, $\tau_S \equiv m_h^2/(4M_S^2)$, and

$$A_0(\tau) \equiv -\frac{1}{\tau^2}[\tau - f(\tau)]. \tag{10}$$

We find $A_0 \simeq 0.37 - 0.34$ for $M_S = 150$–400 GeV, so that an appropriate value of the Higgs-portal coupling is

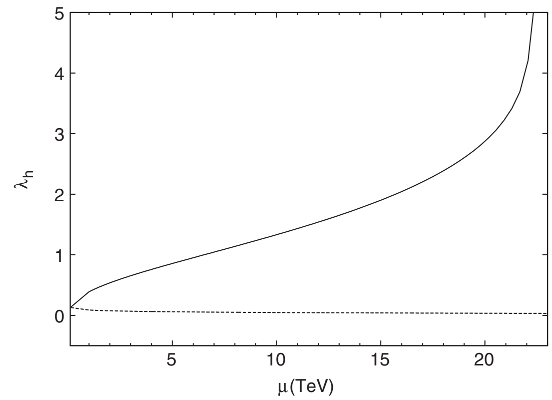

FIG. 1. The running behavior of the IS Higgs quartic coupling λ_h. The solid and dashed lines correspond to λ_h and the SM Higgs quartic coupling, respectively. We fixed the IS Higgs mass $m_h = \sqrt{2\lambda_h}v_h = 125$ GeV and took $\lambda_{hS} = 1.8$ and $\lambda_S = 1.5$. Unlike the SM, the IS Higgs quartic coupling grows up due to a large Higgs-portal coupling λ_{hS} and a small top-Yukawa coupling y_t.

$$\lambda_{hS} \simeq 2.5\text{–}2.7 \times \frac{M_S^2}{vv_h}. \tag{11}$$

As a typical value, we may take $\lambda_{hS} = 1.8$ for $M_S = 200$ GeV and $v_t = 50$ GeV. When $m_{0,S}^2 \approx 0$, i.e., $M_S^2 \approx \lambda_{hS}v_h^2/2$, we obtain $\Gamma(h\to gg) \approx 0.6 \times \Gamma^{\rm SM}(H\to gg)$, independently of the values of λ_{hS}. In order to stabilize the Higgs potential for S at the tree level, the relation $|\lambda_{hS}| < 2\sqrt{\lambda_S\lambda_h}$ is also required.

A comment concerning the IS Higgs quartic coupling λ_h is in order. In the SM, the Higgs mass 125 GeV suggests that the theory is perturbative up to an extremely high energy scale [22]. On the contrary, in the present model, when we take a large Higgs-portal coupling λ_{hS} that reproduces $gg \to h$ correctly, the quartic coupling λ_h will grow because the β-function for λ_h contains the λ_{hS}^2 term. Also, there is no large negative contribution to the β-function for λ_h from the top-Yukawa coupling $y_t \sim 10^{-2}$.

One can demonstrate such a behavior more explicitly by using the renormalization group equations. In Fig. 1, the running of the coupling λ_h is shown. The IS Higgs mass is $m_h = \sqrt{2\lambda_h}v_h$, and we take it to be equal to 125 GeV. Taking a large Higgs-portal coupling $\lambda_{hS} = 1.8$ and the S^4-coupling $\lambda_S = 1.5$, it turns out that the coupling λ_h rapidly grows. Due to the running effects, the naive instability of the scalar potential at the tree level is resolved around the TeV scale in this case. The blowup scale strongly depends on the initial values of λ_{hS} and λ_S. A detailed analysis will be performed elsewhere. Last but not least, we would like to mention that other realizations of the enhancement of the h boson production are also possible.

V. QUARK MASS MATRICES

Let us discuss the structure of the quark mass matrices in the present model. The Yukawa interactions are written by [6,7]

$$-\mathcal{L}_Y = \sum_{i,j} \bar{\psi}_L^{(i)} Y_D^{ij} d_R^{(j)} \Phi_h + \sum_{i,j} \bar{\psi}_L^{(i)} Y_U^{ij} u_R^{(j)} \tilde{\Phi}_h + y_{h_t} \bar{\psi}_L^{(3)} t_R \tilde{\Phi}_{h_t}, \tag{12}$$

with

$$\tilde{\Phi}_h \equiv i\tau_2 \Phi_h^*, \qquad \tilde{\Phi}_{h_t} \equiv i\tau_2 \Phi_{h_t}^*,$$
$$\Phi_{h_t} = \begin{pmatrix} \omega_t^+ \\ \frac{1}{\sqrt{2}}(v_t + h_t + iz_t) \end{pmatrix}, \tag{13}$$
$$\langle \Phi_h \rangle = \begin{pmatrix} 0 \\ \frac{v_h}{\sqrt{2}} \end{pmatrix}, \qquad \langle \Phi_{h_t} \rangle = \begin{pmatrix} 0 \\ \frac{v_t}{\sqrt{2}} \end{pmatrix},$$

$$Y_D \equiv \frac{\sqrt{2}}{v_h} M_D, \qquad Y_U \equiv \frac{\sqrt{2}}{v_h} M_U, \tag{14}$$

and

$$M_D = \begin{pmatrix} m_0^{(1)} & \xi_{12} m_0^{(1)} & \xi_{13} m_0^{(1)} \\ \xi_{21} m_0^{(1)} & m_0^{(2)} + \delta \cdot m_b & \xi_{23} m_0^{(2)} \\ \xi_{31} m_0^{(1)} & \xi_{32} m_0^{(2)} & m_0^{(3)} \end{pmatrix},$$
$$M_U = \begin{pmatrix} \eta_{11} m_0^{(1)} & \eta_{12} m_0^{(1)} & \eta_{13} m_0^{(1)} \\ \eta_{21} m_0^{(1)} & m_0^{(2)} + \delta \cdot m_t & \eta_{23} m_0^{(2)} \\ \eta_{31} m_0^{(1)} & \eta_{32} m_0^{(2)} & m_0^{(3)} \end{pmatrix}, \tag{15}$$

where $\psi_L^{(i)}$ denotes the weak doublet quarks from the i-th family, and $u_R^{(i)}$ and $d_R^{(i)}$ represent the right-handed up- and down-type quarks. The top-Higgs part is responsible for the top mass, $m_t \simeq y_{h_t} \frac{v_t}{\sqrt{2}}$. The IS masses $m_0^{(i)}$ are the same mass scales as the down-type quarks, say, $m_0^{(3)} \sim 1$ GeV, $m_0^{(2)} \sim 100$ MeV, and $m_0^{(1)} \sim 1$ MeV. The common one-loop factor $\delta \sim 1/100$ yields the correct mass hierarchy between m_s and m_c via the hierarchy between m_b and m_t. Also, the off-diagonal coefficients are assumed to be $\xi_{ij}, \eta_{ij} \sim \mathcal{O}(1)$, with some dynamical mechanism. (We kept η_{11} in the up sector for generality.) The CKM matrix is approximately determined by the down-type quark mass matrix [7],

$$V_{\rm CKM} \approx \begin{pmatrix} 1 - \frac{|\xi_{12}|^2}{2}\left(\frac{m_0^{(1)}}{m_0^{(2)}}\right)^2 & \xi_{12}\frac{m_0^{(1)}}{m_0^{(2)}} & \xi_{13}\frac{m_0^{(1)}}{m_0^{(3)}} \\ -\xi_{12}^*\frac{m_0^{(1)}}{m_0^{(2)}} & 1 - \frac{|\xi_{12}|^2}{2}\left(\frac{m_0^{(1)}}{m_0^{(2)}}\right)^2 & \xi_{23}\frac{m_0^{(2)}}{m_0^{(3)}} \\ -(\xi_{13}^* - \xi_{12}^*\xi_{23}^*)\frac{m_0^{(1)}}{m_0^{(3)}} & -\xi_{23}^*\frac{m_0^{(2)}}{m_0^{(3)}} & 1 \end{pmatrix}. \tag{16}$$

We can then reproduce the CKM matrix, basically. For example, with the inputs, $m_0^{(1)} = 10$ MeV, $m_0^{(2)} = 68$ MeV, $m_0^{(3)} = 4.2$ GeV, $m_t = 173.5$ GeV, $\delta = 7 \times 10^{-3}$, $\xi_{12} = \xi_{21} = \eta_{12} = \eta_{21} = 2.0$, $\xi_{13} = \xi_{31} = \eta_{13} = \eta_{31} = 1.6$, $\xi_{23} = \xi_{32} = \eta_{23} = \eta_{32} = -2.5$, $\eta_{11} = \frac{1}{4}$, we obtain $m_d = 4.9$ MeV, $m_s = 95$ MeV, $m_b = 4.2$ GeV, $m_u = 2.2$ MeV, $m_c = 1.3$ GeV, $|V_{ud}| \simeq |V_{cs}| = 0.975$, $|V_{tb}| \simeq 1$, $|V_{us}| \simeq |V_{cd}| = 0.22$, $|V_{cb}| = 0.041$, $|V_{ts}| = 0.039$, $|V_{ub}| = 0.0042$, $|V_{td}| = 0.013$. These values fairly agree with the PDG ones [8].

As was emphasized above in Sec. II, because of the subcriticality dynamics in this class of models, the extra bosons from the top Higgs doublet Φ_{h_t} are heavy, say, $\mathcal{O}(1$ TeV). Thus, their one-loop contributions to the B^0-$\bar{B}^0$ mixing, $b \to s\gamma$ and $Z \to b\bar{b}$ are suppressed. A tree FCN current term also appears in the up sector, so that the D^0-$\bar{D}^0$ mixing is potentially dangerous. However, because the FCN current coupling Y_{t-c-h_t} is found to be tiny, $Y_{t-c-h_t} \sim \frac{m_t}{v_t} \times \frac{m_u}{m_t}\frac{m_c}{m_t} = 10^{-6}$–$10^{-7}$, this does not cause any troubles.

VI. CONCLUSION

The model with an IS Higgs boson yields not only an explanation of the ATLAS and CMS data, including the enhanced diphoton Higgs decay rate, but also makes several predictions. The most important of them is that the value of the top-Yukawa coupling h-t-$\bar{t}$ should be close to the bottom-Yukawa one. Another prediction relates to the decay mode $h \to Z\gamma$, which unlike $h \to \gamma\gamma$ is enhanced only slightly, $\Gamma^{\rm IS}(h \to Z\gamma) = 1.07 \times \Gamma^{\rm SM}(H \to Z\gamma)$. Last but not least, the LHC might potentially discover the top-Higgs resonance h_t, if lucky.

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