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LHC Higgs production and decay in the T' model

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At $\sqrt{s} = 7$ TeV, the standard model (SM) needs at least 10 $(\text{fb})^{-1}$ integrated luminosity at LHC to make a definitive discovery of the Higgs boson. Using binary tetrahedral (T') discrete flavor symmetry, we discuss how the decay of the lightest T' Higgs into $\gamma\gamma$ can be effectively enhanced and dominate over its decay into $b\bar{b}$. Since the two-photon final state allows for a clean reconstruction, a decisive Higgs discovery may be possible at 7 TeV with the integrated luminosity only of ~ 1 $(\text{fb})^{-1}$.

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I. INTRODUCTION

The standard model (SM) of particle physics is a well-tested theory which successfully predicts the strong and electroweak interactions of elementary particles. While its predictive power is impressive, it has limitations. First, in its minimal form, neutrinos have been introduced as massless particles. However, a wealth of experimental data have confirmed that neutrinos are massive and that flavors can mix. Undoubtedly, neutrino mixing is the first indisputable physics beyond the minimal SM. Moreover, in order to specify the SM and make predictions, we need approximately 28 free parameters, including the gauge couplings, quark and lepton masses, mixing angles and possible CP-violating phases, etc.

Grand unification theories (GUT), with or without supersymmetry (SUSY), have been invoked to explain the origin of these free parameters or the relationship between them [1–3]. These models usually focus on reducing the number of parameters in the gauge sector and either the quark or lepton sector, but not both.

A notable alternative to GUTs are models constructed with discrete flavor symmetry. Here, we will focus on the binary tetrahedral group T', which provides calculability to both quark and lepton sectors [4–6]. This model relates quarks and leptons through the T' symmetry, whose irreducible representations are three singlets, three doublets, and a triplet. Since different quark families are assigned to T' singlets and doublets, mass hierarchies in the quark sector appear naturally in the quark sector. Also, the fact that all the $SU_L(2)$ lepton doublets are assigned to a T' triplet is to some extent a unification of the lepton sector. The renormalizable T' model has led to successful predictions of the tribimaximal neutrino-mixing matrix as well as the Cabibbo angle [5,6]. Recently, it has been shown that the discrepancy between the SM prediction and experimental value of muon $g - 2$ factor can be easily accommodated in this model [7]. More details about the T' model, its variants, and other related models can be found in [8]. The success of the renormalizable T' model inspires us to ask if it can be tested at the LHC. In this article, we study Higgs production and decay in the T' model at the LHC.

SM Higgs production and decays have been studied in considerable detail [9]. For instance, due to the high gluon luminosity, gluon-gluon fusion $gg \to h$ is the dominant Higgs production mechanism at the LHC for Higgs masses up to $M_h \sim 1$ TeV. This is about an order of magnitude larger than the next-most important production process, $q\bar{q}' \to hW^{\pm}$. The gluon coupling to Higgs is mediated by the triangular quark loops, and the process is dominated by the top and bottom loops. For $M_h \lesssim 160$ GeV, the branching ratio for the decay process $h \to b\bar{b}$ dominates over all other decay processes, such as $h \to gg$, $h \to WW^*$, $h \to ZZ^*$, and $h \to \gamma\gamma$. In this Higgs mass regime, the production of all other fermion pairs is relatively suppressed compared to $b\bar{b}$ because they are either produced by a relatively lower branching ratio through Higgs decay or by mixing. For $M_h > 160$ GeV, the branching ratio for the decay process $h \to WW$ will take over and dominate.

Given the current limited luminosity of LHC at $\sqrt{s} = 7$ TeV, the more relevant mass range for the SM Higgs would be $M_h \lesssim 160$ GeV. Because of the large QCD background, it is difficult to confirm the processes $h \to b\bar{b}$ and $h \to gg$. While $h \to WW^*$ and $h \to ZZ^*$ have relatively higher rates, the analysis is complicated by escaping neutrinos. As a consequence, in the regime $M_h \lesssim 160$ GeV, the cleanest signal would be $h \to \gamma\gamma$ despite its tiny rate.

In this article, we focus on the decays of the lightest T' Higgs with mass less than 160 GeV. We compare the decay rates of T' Higgses with those of the SM. In particular, we will show that in the fermiophobic limit, the decay of the lightest T' Higgs into $\gamma\gamma$ is effectively enhanced and dominates over its decay into $b\bar{b}$. Since the high p_{T} two-photon final state allows a clean reconstruction, a decisive

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Higgs discovery may be possible in this limit, even at $\sqrt{s} = 7$ TeV with the integrated luminosity of only $\sim 1\ (\mathrm{fb})^{-1}$. As a bonus, the lightest T' Higgs can be unambiguously distinguished from the SM Higgs.

II. THE T' MODEL AND LIGHTEST HIGGS BOSON

We start with a brief review of the simplified model proposed in [6] based on the global symmetry $(T' \times Z_2)$. In particular, we will ignore the lepton sector, which will not be relevant to our study in this article. Interested readers can refer to [4–6] for more details.

In the T' model, left-handed quark doublets $(t, b)_L$, $(c, d)_L$, $(u, d)_L$ are assigned under this global symmetry as

$$\begin{array}{lll} \begin{pmatrix} t \\ b \end{pmatrix}_L & Q_L & (\mathbf{1_1}, +1) \\ \left.\begin{array}{l} \begin{pmatrix} c \\ c \end{pmatrix}_L \\ \begin{pmatrix} u \\ d \end{pmatrix}_L \end{array}\right\} & \mathcal{Q}_L & (\mathbf{2_1}, +1) \end{array}, \tag{1}$$

and the six right-handed quarks as

$$\begin{array}{ll} t_R & (\mathbf{1_1}, +1) \\ b_R & (\mathbf{1_2}, -1) \\ \left.\begin{array}{l} c_R \\ u_R \end{array}\right\} \mathcal{C}_R & (\mathbf{2_3}, -1) \\ \left.\begin{array}{l} s_R \\ d_R \end{array}\right\} \mathcal{S}_R & (\mathbf{2_2}, +1). \end{array} \tag{2}$$

The quark-Yukawa sector in the model is thus given as [6]

$$\mathcal{L}_Y^q = Y_t(\{Q_L\}_{\mathbf{1_1}}\{t_R\}_{\mathbf{1_1}} H_{\mathbf{1_1}}) + Y_b(\{Q_L\}_{\mathbf{1_1}}\{b_R\}_{\mathbf{1_2}} H_{\mathbf{1_3}}) + Y_{\mathcal{C}}(\{\mathcal{Q}_L\}_{\mathbf{2_1}}\{\mathcal{C}_R\}_{\mathbf{2_3}} H'_{\mathbf{3}}) + Y_{\mathcal{S}}(\{\mathcal{Q}_L\}_{\mathbf{2_1}}\{\mathcal{S}_R\}_{\mathbf{2_2}} H_{\mathbf{3}}) + \text{H.c.}. \tag{3}$$

A. T'-Higgs couplings

We focus on flavor-diagonal interactions and fermion couplings to a set of neutral T'-Higgs bosons, $\{H_r^{(i)}\}$, where $r = (1_1, 1_3, 3, 3')$ denotes the T'-irreducible representation and i denotes components in the T'-multiplet.[1] Let the Higgs vacuum expectation values (VEVs) be $\langle H_r^{(i)} \rangle = v_r^{(i)}/\sqrt{2}$. Expanding fields in terms of the Clebsch-Gordan coefficients [6], we have

$$\begin{aligned} \mathcal{L}_Y^q\big|^{\text{neutral}}_{\text{flavor-diagonal}} &= Y_t \bar{t} t \frac{(H_{1_1} + v_{1_1})}{\sqrt{2}} + Y_b \bar{b} b \frac{(H_{1_3} + v_{1_3})}{\sqrt{2}} \\ &\quad - \frac{Y_{\mathcal{C}}}{\sqrt{6}} \bar{c} c \frac{(H_{3'}^{(1)} + v_{3'}^{(1)})}{\sqrt{2}} + \frac{Y_{\mathcal{S}}}{\sqrt{3}} \bar{s} s \frac{(H_3^{(1)} + v_3^{(1)})}{\sqrt{2}} \\ &\quad + \frac{Y_{\mathcal{S}}}{\sqrt{6}} \bar{d} d \frac{(H_3^{(1)} + v_3^{(1)})}{\sqrt{2}} + \sqrt{\frac{2}{3}} Y_{\mathcal{C}} \bar{u} u \frac{(H_{3'}^{(2)} + v_{3'}^{(2)})}{\sqrt{2}}. \end{aligned} \tag{4}$$

The relevant Yukawa couplings and fermion masses are thus read off:

$$g_{H_{1_1} tt} = Y_t, \qquad m_t = \frac{Y_t}{\sqrt{2}} v_{1_1}, \tag{5}$$

$$g_{H_{1_3} bb} = Y_b, \qquad m_b = \frac{Y_b}{\sqrt{2}} v_{1_3}, \tag{6}$$

$$g_{H_{3'}^{(1)} cc} = \frac{Y_{\mathcal{C}}}{\sqrt{6}}, \qquad m_c = \frac{Y_{\mathcal{C}}}{2\sqrt{3}} v_{3'}^{(1)}, \tag{7}$$

$$g_{H_3^{(1)} ss} = \frac{Y_{\mathcal{S}}}{\sqrt{3}}, \qquad m_s = \frac{Y_{\mathcal{S}}}{\sqrt{6}} v_3^{(1)}, \tag{8}$$

$$g_{H_3^{(1)} dd} = \frac{Y_{\mathcal{S}}}{\sqrt{6}}, \qquad m_d = \frac{Y_{\mathcal{S}}}{2\sqrt{3}} v_3^{(1)}, \tag{9}$$

$$g_{H_{3'}^{(2)} uu} = \sqrt{\frac{2}{3}} Y_{\mathcal{C}}, \qquad m_u = \frac{Y_{\mathcal{C}}}{\sqrt{3}} v_{3'}^{(2)}. \tag{10}$$

Fixing fermion masses to be the same as those in the SM, namely $m_f = \frac{g_{hff}^{\text{SM}}}{\sqrt{2}} v_{\text{EW}}$, we may express the T'-Yukawa couplings comparing those in the SM to get

$$g_{H_{1_1} tt} = \left(\frac{v_{\text{EW}}}{v_{1_1}}\right) g_{htt}^{\text{SM}}, \tag{11}$$

$$g_{H_{1_3} bb} = \left(\frac{v_{\text{EW}}}{v_{1_3}}\right) g_{hbb}^{\text{SM}}, \tag{12}$$

$$g_{H_{3'}^{(1)} cc} = \left(\frac{v_{\text{EW}}}{v_{3'}^{(1)}}\right) g_{hcc}^{\text{SM}}, \tag{13}$$

$$g_{H_3^{(1)} ss} = \left(\frac{v_{\text{EW}}}{v_3^{(1)}}\right) g_{hss}^{\text{SM}}, \tag{14}$$

$$g_{H_3^{(1)} dd} = \left(\frac{v_{\text{EW}}}{v_3^{(1)}}\right) g_{hdd}^{\text{SM}}, \tag{15}$$

$$g_{H_{3'}^{(2)} uu} = \left(\frac{v_{\text{EW}}}{v_{3'}^{(2)}}\right) g_{huu}^{\text{SM}}. \tag{16}$$

We next turn to the gauge-Higgs sector, $\mathcal{L}_{\text{GH}} = \sum_{r,i} |D_\mu H_r^{(i)}|^2$, where all the T'-Higgs fields couple to the electroweak gauge bosons. The W and Z boson masses are thus expressed in terms of the T'-Higgs VEVs as follows:

[1] As was studied in the model building of Ref. [8], if we go beyond the minimal T' model to incorporate mixing with the third generation of quarks, we may encounter flavor-changing neutral current (FCNC) problems because the Higgs bosons can have off-diagonal flavor couplings, unlike the SM Higgs. Then, the size of T' Yukawa couplings to the mass-eigenstate Higgs bosons, namely $Y_f^2 (a_n^f)^2 / M_{H_n}^2 \propto (a_n^f/R_f)^2 (g_{hff}^{\text{SM}})^2 M_{H_n}^2$, would be constrained by the FCNC issue, which may give further constraints on the parameters a_n^f and R_f. More on this issue is beyond scope of the present article and is to be pursued in detail in the future.

$$M_W^2 = \frac{g_W^2}{4}\sum_{r,i}(v_r^{(i)})^2, \qquad M_Z^2 = \frac{M_W^2}{c_W^2}, \tag{17}$$

where g_W is the $SU(2)_W$ gauge coupling and $c_W = \frac{g_W}{\sqrt{g_W^2+g_Y^2}}$ with g_Y being $U(1)_Y$ gauge coupling.[2] We fix the W and Z boson masses to be those in the SM. This can be achieved by identifying the electroweak scale $v_{\rm EW}$ as

$$v_{\rm EW}^2 = \sum_{r,i}(v_r^{(i)})^2, \tag{18}$$

so that we have $M_W^2 = g_W^2 v_{\rm EW}^2/4$.

The T'-Higgs couplings to WW and ZZ read

$$\mathcal{L}_{H_r^{(i)}VV} = g_{H_r^{(i)}WW}H_r^{(i)}W_\mu^+W^{\mu-} + \frac{1}{2}g_{H_r^{(i)}ZZ}H_r^{(i)}Z_\mu Z^\mu, \tag{19}$$

where

$$g_{H_r^{(i)}WW} = g_W^2 v_r^{(i)} = \left(\frac{v_r^{(i)}}{v_{\rm EW}}\right)g_{hWW}^{\rm SM}, \tag{20}$$

$$g_{H_r^{(i)}ZZ} = \frac{g_W^2}{c_W^2}v_r^{(i)} = \left(\frac{v_r^{(i)}}{v_{\rm EW}}\right)g_{hZZ}^{\rm SM}, \tag{21}$$

with $g_{hVV}^{\rm SM}$ ($V = W, Z$) being the corresponding coupling to the Higgs boson in the SM,

$$g_{hVV}^{\rm SM} = 4\frac{M_V^2}{v_{\rm EW}}. \tag{22}$$

B. The lightest Higgs boson and its couplings

Electroweak interactions mix T'-Higgs doublets. Given an explicit form of Higgs potential, we can solve such a mixing to get a set of mass eigenstates $\{H_n\}$ with their eigenfunctions, a_n^r. Without knowing the explicit expression of Higgs potential, in general, we may write

$$H_r^{(i)} = \sum_n a_n^{(r,i)}H_n, \tag{23}$$

where the expansion coefficients $a_n^{(r,i)}$ form an orthonormal complete set,

$$\sum_{r,i}a_n^{(r,i)}a_m^{(r,i)} = \delta_{nm}, \tag{24}$$

which followed from the normalization condition of the kinetic terms of $\{H_n\}$.

Assuming a mass hierarchy for the Higgs bosons, $M_{H_0} < M_{H_1} < \cdots$, we can identify the lightest Higgs boson as H_0 with mass, say, $\lesssim m_t \simeq 172$ GeV. Hereafter, we shall confine ourselves to the phenomenology of this H_0.

[2]Note that we have $\rho = 1$ at tree level as in the SM. This is because T' symmetry commutes with the electroweak symmetry as well as the custodial symmetry.

It is convenient to introduce a ratio,

$$R_r^{(i)} = \frac{v_r^{(i)}}{v_{\rm EW}}, \tag{25}$$

which satisfies

$$\sum_{r,i}(R_r^{(i)})^2 = 1. \tag{26}$$

From Eqs. (11)–(16), (20), and (21), we then obtain the H_0 couplings to fermions,

$$g_{H_0tt} = \left(\frac{a_0^{1_1}}{R_{1_1}}\right)g_{htt}^{\rm SM}, \tag{27}$$

$$g_{H_0bb} = \left(\frac{a_0^{1_3}}{R_{1_3}}\right)g_{hbb}^{\rm SM}, \tag{28}$$

$$g_{H_0cc} = \left(\frac{a_0^{(3',1)}}{R_{3'}^{(1)}}\right)g_{hcc}^{\rm SM}, \tag{29}$$

$$g_{H_0ss} = \left(\frac{a_0^{(3,1)}}{R_3^{(1)}}\right)g_{hss}^{\rm SM}, \tag{30}$$

$$g_{H_0dd} = \left(\frac{a_0^{(3,1)}}{R_3^{(1)}}\right)g_{hdd}^{\rm SM}, \tag{31}$$

$$g_{H_0uu} = \left(\frac{a_0^{(3',2)}}{R_{3'}^{(2)}}\right)g_{huu}^{\rm SM}, \tag{32}$$

and gauge bosons

$$g_{H_0VV} = \sum_{r,i}(a_0^{(r,i)}R_r^{(i)})g_{hVV}^{\rm SM}, \tag{33}$$

where $V = W, Z$.

III. LHC HIGGS DECAY AND PRODUCTION

In this section, we study the decay modes of a light Higgs boson with mass in the range 109 GeV $\lesssim M_{H_0} \lesssim (2M_W \simeq)160$ GeV.[3] In this mass range, $h \to b\bar{b}$ and $h \to gg$ are dominant modes in the SM, where the top and bottom loops give the significant effect on the $h \to gg$ mode. Here, we shall focus on the top and bottom contributions to decay modes of the T'-Higgs boson H_0. The relevant formulas for its partial decay widths of H_0 are given in Appendix A.

From Eqs. (27)–(33), we see that the difference between the SM and the T' model is handled by two kinds of parameters, $a_0^{(r,i)}$ and $R_r^{(i)}$. Since we are interested in the

[3]The lower bound comes from the exclusion limit of the direct search of $h \to \gamma\gamma$ at the LEP II [10].

top and bottom contributions, we may take $a_0^{(r,i)} = 0$ except $a_0^{1_1}$ and $a_0^{1_3}$ in Eqs. (27)–(32) and keep only R_{1_1} and R_{1_3} nonzero in Eq. (33), so that all the Yukawa couplings other than those of top and bottom vanish and the H_0-V-V coupling is saturated only by H_{1_1} and H_{1_3} in the sum. Note that Eqs. (24) and (26) then constrain the remaining parameters:

$$(a_0^{1_1})^2 + (a_0^{1_3})^2 = 1, \qquad 0 \le a_0^{1_1} \le 1, \qquad 0 \le a_0^{1_3} \le 1,$$
$$(R_{1_1})^2 + (R_{1_3})^2 = 1, \qquad 0 \le R_{1_1} \le 1, \qquad 0 \le R_{1_3} \le 1. \tag{34}$$

From Eqs. (27), (28), (33), and (34), one can see that the SM limit is given by

$$a_0^{1_3} \to R_{1_3} \to 0, \tag{35}$$

in such a way that $g_{H_0tt} \to g_{htt}^{\rm SM}$, $g_{H_0bb} \to g_{hbb}^{\rm SM}$, and $g_{H_0VV} \to g_{hVV}^{\rm SM}$. On the other hand, a fermiophobic limit can be taken as

$$a_0^{1_3} \to 0, \tag{36}$$

in a sense that the bottom Yukawa coupling goes to zero and $H_0 \to b\bar{b}$ mode gets highly suppressed to be zero, while the top Yukawa coupling remains nonzero.

In Figs. 1–3, we show the branching fraction of the lightest T' Higgs decay (left panels) and the ratio to that of the SM Higgs (right panels). As a sample, we have taken $a_0^{1_3} = 2/3, 1/3, 0$ with $R_{1_3} = 0.5$ fixed, which monitors the interpolation between the SM case and the fermiophobic

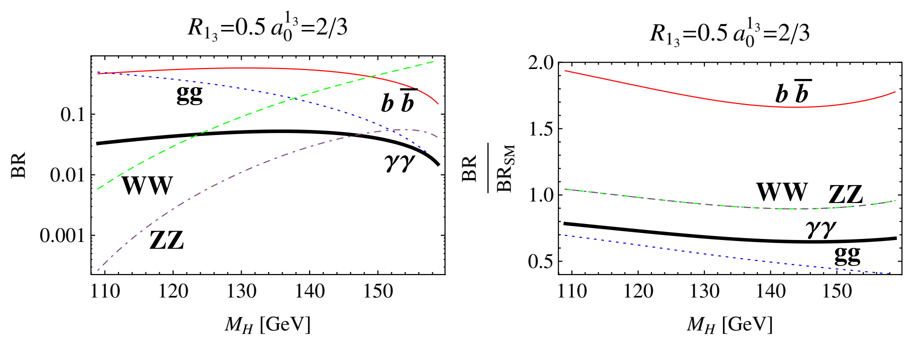

FIG. 1 (color online). Branching fraction of the lightest T'-Higgs boson with $R_0^{1_3} = 0.5$ and $a_0^{1_3} = 2/3$.

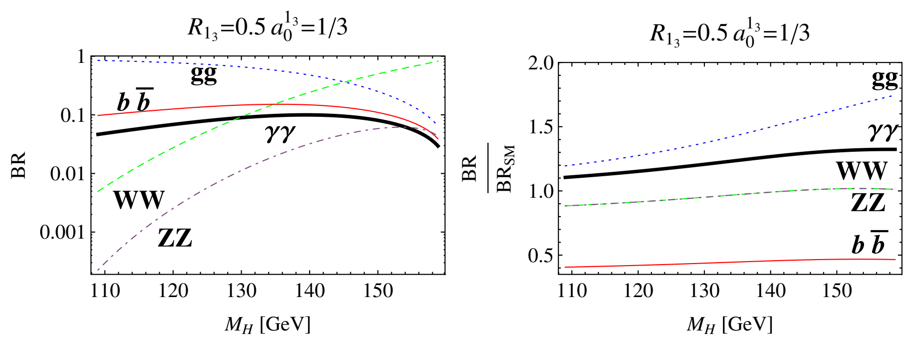

FIG. 2 (color online). Branching fraction of the lightest T'-Higgs boson with $R_0^{1_3} = 0.5$ and $a_0^{1_3} = 1/3$.

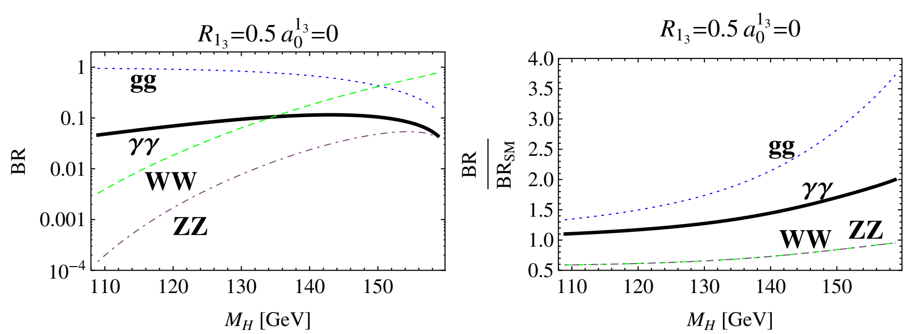

FIG. 3 (color online). Branching fraction of the lightest T'-Higgs boson with $R_0^{1_3} = 0.5$ and $a_0^{1_3} = 0$, in which case the Higgs boson does not couple to $b\bar{b}$ since the Yukawa coupling vanishes.

case. Figures. 1–3 imply that as $a_0^{1_3}$ approaches the fermiophobic limit $a_0^{1_3} \to 0$, $H_0 \to \gamma\gamma$ becomes dominant in contrast to the case of the SM Higgs, in which $h \to b\bar{b}$ is dominant. Note that the fermiophobicity does not affect WW and ZZ decay modes so much (See Fig. 3), since $g_{H_0VV}/g_{hVV}^{\rm SM} = (\sqrt{1-(a_0^{1_3})^2}\sqrt{1-R_{1_3}^2} + a_0^{1_3}R_{1_3}) \to \sqrt{1-R_{1_3}^2}$ when $a_0^{1_3} \to 0$.

In Figs. 4–6, we show contour plots of $\mathrm{Br}(H_0 \to b\bar{b})$, $\mathrm{Br}(H_0 \to \gamma\gamma)$, and $\mathrm{Br}(H_0 \to gg)$ for $M_H = 120$ GeV in the entire region of the parameter space $(a_0^{1_3}, R_{1_3})$, comparing with those of the SM Higgs. It is interesting to note from Figs. 4–6 that $H_0 \to b\bar{b}$ mode is necessarily suppressed when $H_0 \to \gamma\gamma$ mode is enhanced, while $H_0 \to gg$ mode is enhanced at the same time, which is due to the remaining sizable top loop contribution: One cannot make both top and bottom quarks decoupled simultaneously because of the constraint (34).

Finally, let us briefly discuss Higgs production. In particular, in the true fermiophobic limit, which is realized by taking $a_0^{1_3} \to 0$, we find $a_0^{1_1} \to 1$. This implies that the gluon-gluon fusion through the top-triangular loop will be the dominant production process for the fermiophobic T'-Higgs, which is the same (within a percent) as in the SM. Even though the dominant Higgs production process is the same, the fermiophobic T'-Higgs can be unambiguously distinguished from the SM Higgs at the LHC because $H_0 \to \gamma\gamma$ dominates over $H_0 \to b\bar{b}$ in the range 109 GeV $\lesssim M_{H_0} \lesssim$ 160 GeV. Since $H_0 \to \gamma\gamma$ allows for a clean reconstruction, a decisive Higgs discovery may be possible even at $\sqrt{s} = 7$ TeV with the integrated luminosity of only $\sim 1(\text{ fb})^{-1}$.

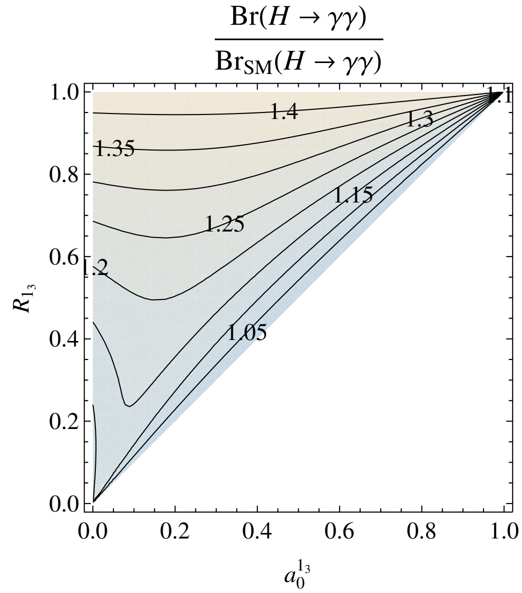

FIG. 5 (color online). Contour plot of $\mathrm{Br}(H_0 \to \gamma\gamma)/\mathrm{Br}_{\rm SM}(H_0 \to \gamma\gamma)$ on the $(a_0^{1_3}, R_{1_3})$ plane for $M_H = 120$ GeV. The contour points are restricted to a region where the value is larger than 1. The limit $a_0^{1_3} \to R_{1_3} \to 0$ along the line $R_{1_3} = a_0^{1_3}$ corresponds to the SM case.

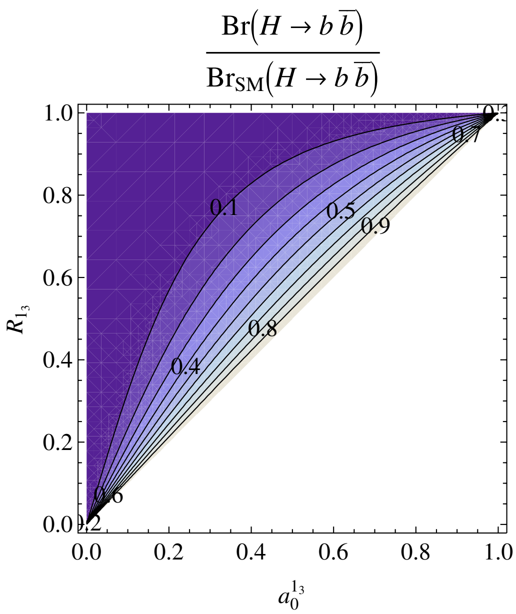

FIG. 4 (color online). Contour plot of $\mathrm{Br}(H_0 \to b\bar{b})/\mathrm{Br}_{\rm SM}(H_0 \to b\bar{b})$ on the $(a_0^{1_3}, R_{1_3})$ plane for $M_H = 120$ GeV. The contour points are restricted to a region where the value is less than 1. The limit $a_0^{1_3} \to R_{1_3} \to 0$ along the line $R_{1_3} = a_0^{1_3}$ corresponds to the SM case.

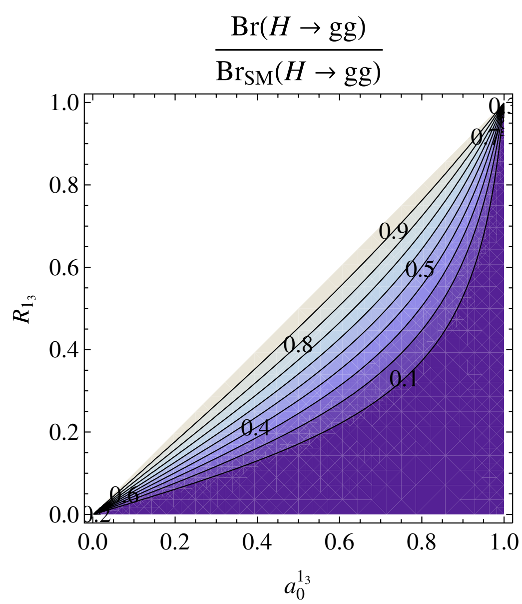

FIG. 6 (color online). Contour plot of $\mathrm{Br}(H_0 \to gg)/\mathrm{Br}_{\rm SM}(H_0 \to gg)$ on the $(a_0^{1_3}, R_{1_3})$ plane for $M_H = 120$ GeV. The contour points are restricted to a region where the value is less than 1. The limit $a_0^{1_3} \to R_{1_3} \to 0$ along the line $R_{1_3} = a_0^{1_3}$ corresponds to the SM case.

IV. DISCUSSION

This article may be taken as a warning to experimentalists that the properties of the lightest Higgs boson can readily depart very significantly from the predictions of the minimal SM with only one Higgs doublet, and with its 28 parameters unconstrained by any further theoretical input.

We have studied a model with a $(T' \times Z_2)$ flavor symmetry which commutes with the SM gauge group, and which leads to agreement with the mixing matrices for neutrinos and quarks. It necessarily changes the couplings of the lightest Higgs to the quarks and leptons, which are no longer simply proportional to the fermion masses. This aspect of the SM is its most fragile prediction.

A similar, but different, illustration of this fragility is provided by the variant axion model [11], instead motivated by solution of the strong *CP* problem. In both cases, the delimiting of the SM parameters change in the Yukawa sector.

In the present case, the T' flavor symmetry can give rise to optimism that the discovery of the Higgs may be expedited because the product of the production cross-section and the decay branching ratio is enhanced. As can be seen from Figs. 3 and 5, the decay $H \to \gamma\gamma$ is generically larger even for $M_H = 120$ GeV and becomes more so at larger Higgs mass. Even with $\sqrt{s} = 7$ TeV and 1 $(\mathrm{fb})^{-1}$, the LHC could make a Higgs discovery.

Many aspects of the SM have been confirmed to high accuracy. These checks are principally for the gauge sector, which has a significant geometrical underpinning and hence also has uniqueness. The Yukawa sector, where most of the 28 free parameters lie, does not have a geometrical interpretation. The objective of the flavor symmetry is to supply an explanation of some of the parameters, and it is therefore interesting to explore other predictions for production and decay of Higgs at the LHC.

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APPENDIX A: FORMULAS FOR HIGGS DECAY WIDTHS

In this appendix, we shall present formulas of decay widths relevant to the H_0 decay modes.

1. $H_0 \to q\bar{q}$ mode

In the SM, the $h \to q\bar{q}$ decay width is calculated at the leading order of perturbation to be

$$\Gamma^{\rm SM}[h \to q\bar{q}] = \frac{N_c M_h}{16\pi}(g^{\rm SM}_{hqq})^2\left(1 - \frac{4m_q^2}{M_h^2}\right)^{3/2}. \quad \text{(A1)}$$

To get the corresponding formula for $H_0 \to q\bar{q}$, all we need to do is replace the Yukawa coupling and the Higgs boson mass with the appropriate ones. Thus, we have

$$\Gamma^{T'}[H_0 \to q\bar{q}] = \frac{N_c M_{H_0}}{16\pi}(g_{H_0qq})^2\left(1 - \frac{4m_q^2}{M_{H_0}^2}\right)^{3/2}. \quad \text{(A2)}$$

2. $H_0 \to gg$ mode

In the SM, we compute the $h \to gg$ decay width at the leading order of perturbation to get

$$\Gamma^{\rm SM}[h \to gg] = \frac{N_c^2\alpha_S^2 M_h^3}{576\pi^3}\left|\sum_q \frac{g^{\rm SM}_{hqq}}{m_q}(1 + (1-\tau_q)f(\tau_q))\tau_q\right|^2, \quad \text{(A3)}$$

where $\tau_q = 4m_q^2/M_h^2$ and defined [9]

$$f(\tau) = \begin{cases} (\sin^{-1}\frac{1}{\sqrt{\tau}})^2 & \tau \geq 1 \\ -\frac{1}{4}(\log(\frac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}}) - i\pi))^2 & \tau < 1 \end{cases}.$$

Replacing $g^{\rm SM}_{hqq}$ with g_{H_0qq} and M_h with M_{H_0}, we have

$$\Gamma^{T'}[H_0 \to gg] = \frac{N_c^2\alpha_S^2 M_{H_0}^3}{576\pi^3}\left|\sum_q \frac{g_{H_0qq}}{m_q}(1 + (1-\tau_q)f(\tau_q))\tau_q\right|^2, \quad \text{(A4)}$$

where $\tau_q = 4m_q^2/M_{H_0}^2$.

3. $H_0 \to \gamma\gamma$ mode

In the SM, the leading contribution to the $h \to \gamma\gamma$ decay width is calculated to be

$$\Gamma^{\rm SM}[h \to \gamma\gamma] = \frac{\alpha^2 M_h^3}{256\pi^3}\left|N_c\sum_q \frac{g^{\rm SM}_{hqq}}{\sqrt{2}m_q}Q_q^2 A_q^h(\tau_q) + \frac{g^{\rm SM}_{hWW}}{4M_W^2}A_W^h(\tau_W)\right|^2, \quad \text{(A5)}$$

where we neglected contributions from lepton-triangle loops, and defined [9]

$$A_q^h(\tau) = 2\tau[1 + (1-\tau)f(\tau)], \quad \text{(A6)}$$

$$A_W^h(\tau) = -[2 + 3\tau + 3\tau(2-\tau)f(\tau)] \quad \text{(A7)}$$

with $\tau_i = 4m_i^2/M_h^2$. Replacing couplings and masses with the appropriate ones, we get

$$\Gamma^{T'}[H_0 \to \gamma\gamma] = \frac{\alpha^2 M_{H_0}^3}{256\pi^3} \left| N_c \sum_q \frac{g_{H_0qq}}{\sqrt{2}m_q} Q_q^2 A_q^{H_0}(\tau_q) + \frac{g_{H_0WW}}{4M_W^2} A_W^{H_0}(\tau_W) \right|^2, \tag{A8}$$

where $\tau_i = 4m_i^2/M_{H_0}^2$.

4. $H_0 \to VV^*$ mode

In the SM, the leading contribution to the $h \to VV^*$ decay width is calculated to be

$$\Gamma^{\rm SM}[h \to VV^*] = \delta_{V'} \frac{3G_F(g_{hVV}^{\rm SM})^2 M_h}{256\sqrt{2}\pi^3} R\left(\frac{M_V^2}{M_h^2}\right), \tag{A9}$$

where $G_F = \frac{1}{\sqrt{2}v_{\rm EW}^2}$ and we defined [9]

$$\delta_{V'} = \begin{cases} 1 & \text{for } W \\ \frac{7}{12} - \frac{10}{9}s_W^2 + \frac{40}{27}s_W^4 & \text{for } Z \end{cases}, \tag{A10}$$

$$R(x) = \frac{3(1-8x+20x^2)}{\sqrt{4x-1}} \cos^{-1}\left(\frac{3x-1}{2x^{3/2}}\right) - \frac{(1-x)(2-13x+47x^2)}{2x} - \frac{3}{2}(1-6x+4x^2)\log x. \tag{A11}$$

Replacing $g_{hVV}^{\rm SM}$ and M_h with g_{H_0VV} and M_{H_0}, respectively, we obtain

$$\Gamma^{T'}[H_0 \to VV^*] = \delta_{V'} \frac{3G_F(g_{H_0VV})^2 M_{H_0}}{256\sqrt{2}\pi^3} R\left(\frac{M_V^2}{M_{H_0}^2}\right). \tag{A12}$$

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