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Higgs data confronts sequential fourth generation fermions in the Higgs triplet model

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We investigate the effect of introducing a sequential generation of chiral fermions in the Higgs triplet model with nontrivial mixing between the doublet and triplet Higgs. We use the available Large Hadron Collider data for Higgs boson production and decay rates, the constraints on the fourth generation masses, and impose electroweak precision constraints from the S, T and U parameters. Our analysis shows that a standard model (SM)–like Higgs boson state at ∼125 GeV can be accommodated in the Higgs triplet model with four generations, and thus, that four generations survive collider *and* electroweak precision constraints in models beyond SM.

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I. INTRODUCTION

The new Large Hadron Collider (LHC) data from ATLAS [1] and CMS [2] seems to indicate a Higgs boson that may be consistent with the standard model (SM) one. Although the statistics are not very strong yet, and more analyses are needed, the question on whether the SM is the final theory still remains. In particular, the SM fails to explain observed experimental phenomena such as neutrino masses, dark matter, or baryogenesis. One could reasonably ask if the Higgs boson found at the LHC, even if finally shown to be consistent with the SM predictions, could not possibly belong to a more complete theoretical scenario, where the neutral Higgs signatures are consistent with the SM ones. This situation is akin to requiring new physics scenarios to satisfy low energy precision measurements.

To test this hypothesis, we apply it to one of the simple generalizations of the SM, the Higgs triplet model (HTM), to which a sequential generation of chiral fermions is added. The addition of a fourth sequential generation of fermion doublets is a natural extension of the SM (SM4) [3]. The model restricts fourth generation quark masses from being too large to preserve perturbativity [4], and it does not conflict with electroweak precision observables [3] as long as their mass differences are small [5]. Further limits on the fourth generation fermion masses exist from direct searches at collider experiments such as the Large Electron-Positron Collider (LEP) and Tevatron, as well as from the current LHC data.

There are many advantages of introducing an extra family of fermions:

(i) The fermions associated with the fourth generation could trigger dynamical electroweak symmetry breaking [4] without a Higgs boson, and thus deal with the hierarchy problem.
(ii) Gauge couplings unification can in principle be achieved without supersymmetry [6].
(iii) A new family might resolve SM problems in flavor physics, such as the charge-parity (CP) violation in B_s-mixing [7]. While the electroweak precision data constrains the mass splitting between the fourth generation quarks, data from B-meson physics constrain their mixing pattern [8].
(iv) A fourth generation could solve problems related to baryogenesis because an additional quark doublet could contribute to an increase in the amount of of CP violation [9].
(v) A fourth generation extension of the SM would increase the strength of the electroweak phase transition [10].
(vi) A fourth generation neutrino can serve as a candidate for cold dark matter [11], resolving this outstanding problem of the SM.
(vii) New heavy fermions lead to new interesting effects due to their large Yukawa couplings [12].

However, the SM4 scenario is severely constrained by the available data [8]. First, from constraints on the invisible width of the Z boson at the LEP, the number of light neutrinos is $N_\nu = 2.9840 \pm 0.0084$ [13] and thus the fourth family neutrino must be heavier than $M_Z/2$, assuming small mixing with the lighter SM leptons. A heavy charged lepton with a mass $m_{\ell'} < 100\,\text{GeV}$ has also been excluded at LEP2 [13]. The Tevatron and now the LHC have excluded light fourth generation quarks. Direct

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searches have been performed by both the ATLAS and CMS collaborations, with the CMS collaboration putting the strongest bound on the masses of degenerate fourth generation quarks, ruling out $m_{q'} < 685$ GeV at 95% C.L. [14].[1] Updated bounds can also be extracted from an inclusive search done by the CMS collaboration for vectorlike top partners at $\sqrt{s} = 8$ TeV [15]. Unitarity requirements indicate that fourth generation fermions should not be extremely heavy, $m_{q'} < 500$ GeV [17]. This bound is seen not as a limit, but as an indication that near the perturbative unitarity bound strong dynamics take place.

Strong constraints on SM4 can be also obtained from Higgs searches at the Tevatron and the LHC. As the dominant mode for Higgs production at hadron colliders is through loop induced gluon-gluon fusion, the Higgs–gluon-gluon vertex (hgg) is significantly enhanced by new heavy colored fermions of SM4 in the loop, which couple to the Higgs boson proportionally to their mass. The enhancement in the production cross section can be approximated by a factor of $\frac{\sigma(gg\to h)_{\rm SM4}}{\sigma(gg\to h)_{\rm SM}} \approx 9$ [18]. Such a large enhancement would certainly increase the event rates for a Higgs signal at experiments and, therefore, nonobservation of any signal helps in putting strong constraints on the Higgs boson mass in SM4. The CDF and D0 experiments exclude a Higgs boson in this scenario for masses $124\ \text{GeV} < m_h < 286\ \text{GeV}$ by mainly considering the $gg \to h \to WW \to 2\ell 2\nu$ channel [19]. The LHC experiments recently extended this exclusion limit up to $m_h \approx 600$ GeV (at 99% C.L.) by also exploiting the $gg \to h \to ZZ \to 4\ell, 2\ell 2\nu, 2\ell 2j$ search channels [20]. The new data worsens the situation for SM4 [21,22]. In addition, the limits on the low energy phenomenology due to fourth generation fermions in SM4 has also been studied extensively [23,24].

While there have been many extensive studies of the SM4, there are few analyses of beyond standard model scenarios with four generations (see, however, [25–27]). The reason is that the fourth generation typically imposes severe restrictions on the models. In particular, there are difficulties in incorporating a chiral fourth family scenario into any Higgs doublet model, such as the minimal supersymmetric standard model (MSSM) [28]. It was initially shown that due to the large masses for the fourth generation quarks and large Yukawa couplings, there are no values of $\tan\beta = \frac{v_u}{v_d} > 1$ for which the couplings are perturbative to the grand unification scale. One would need to invoke different couplings, such as one Higgs doublet only coupling to the fourth generation [27]. (However, this condition does not apply to vectorlike quarks [29].) However, the MSSM with four generations has received some more attention [30], as it was shown that for $\tan\beta \simeq 1$ the model exhibits a strong first order phase transition [31]. Four generations can be incorporated naturally into a warped spacetime model [32], seen perhaps as a particular example of composite Higgs models where the Higgs boson emerges as a condensate of the fourth generation fermions [33].

Given the serious shortcomings of SM4, we chose to explore the possibility of a fourth generation model in a simple extension of the SM, the Higgs triplet model [34]. This framework, which we have chosen to call HTM4, immediately has two advantages:

(1) Unlike Higgs doublet models, there are no problems with Yukawa couplings arising from the ratio of the two doublet vacuum expectation values (VEVs), and thus some of the problems with perturbativity are softened.
(2) The Higgs triplet model is the simplest scenario to allow for neutrino masses through the type-II seesaw mechanism [35].

Additionally, we will show that, for small mass splittings within the Higgs multiplets and the additional fermion family, the model satisfies precision conditions on the oblique parameters. We explore whether one of the neutral CP-even Higgs with mass ~125 GeV in the HTM4 could be consistent with the Higgs signals at the LHC.

Our paper is organized as follows. In Sec. II we briefly describe the Higgs triplet model with nontrivial mixing and including a fourth generation. We include the oblique corrections and the restrictions imposed on the model in Sec. III. Our analysis of the parameter space, including collider and precision electroweak constraints, is discussed in Sec. IV, while we present our results in Sec. V. We summarize and conclude in Sec. VI. Additional formulas for production and decay of the Higgs boson(s) are provided in Appendixes A and B.

II. THE HIGGS TRIPLET MODEL WITH FOUR GENERATIONS

The scalar sector of the HTM4 is composed of one isospin doublet field Φ with hypercharge $Y_\Phi = 1$ and one triplet field Δ. It is customary to choose the triplet to be a complex field with hypercharge $Y_\Delta = 2$.[2]

The electric charge is defined as $Q = T_{3L} + \frac{Y}{2}$, with T_{3L} as the third component of the $SU(2)_L$ isospin. The scalar fields Φ and Δ can be parametrized as a 1×2 column and a 2×2 matrix, respectively:

[1] In these experimental bounds, assumptions such as $\text{BR}(b' \to tW)$ or $\text{BR}(t' \to bW) = 1$ are made; relaxing them leads to slightly weaker bounds, as discussed in e.g. [16]. Further softening of the constraints happens for nondegenerate choice of masses for the up-type and down-type quarks of the fourth generation.

[2] A real field with hypercharge $Y_\chi = 0$ is also possible:

$$\chi = \begin{pmatrix} \chi^+ \\ \chi^0 \\ \chi^- \end{pmatrix} \quad \text{with} \quad \chi^0 = \frac{1}{\sqrt{2}}(\chi + v_\chi + i\eta).$$

$$\Phi = \begin{bmatrix} \varphi^+ \\ \frac{1}{\sqrt{2}}(\varphi + v_\Phi + i\chi) \end{bmatrix}, \qquad \Delta = \begin{bmatrix} \frac{\Delta^+}{\sqrt{2}} & \Delta^{++} \\ \Delta^0 & -\frac{\Delta^+}{\sqrt{2}} \end{bmatrix} \quad \text{with}$$

$$\Delta^0 = \frac{1}{\sqrt{2}}(\delta + v_\Delta + i\eta), \tag{1}$$

where v_Φ and v_Δ are the VEVs of the doublet Higgs field and the triplet Higgs field, respectively, which satisfy $v^2 \equiv v_\Phi^2 + 2v_\Delta^2 \simeq (246\ \text{GeV})^2$.

The terms in the Lagrangian relevant for Higgs interactions are given by

$$\mathcal{L}_{\text{HTM}} = \mathcal{L}_{\text{kin}} + \mathcal{L}_Y - V(\Phi, \Delta), \tag{2}$$

where $\mathcal{L}_{\text{kin}}$, $\mathcal{L}_Y$ and $V(\Phi, \Delta)$ are the kinetic term, Yukawa interaction and scalar potential, respectively. The kinetic term for the Higgs fields is

$$\mathcal{L}_{\text{kin}\Phi,\Delta} = (D_\mu \Phi)^\dagger (D^\mu \Phi) + \text{Tr}[(D_\mu \Delta)^\dagger (D^\mu \Delta)], \tag{3}$$

where the covariant derivatives are defined as

$$D_\mu \Phi = \left(\partial_\mu + i\frac{g}{2}\tau^a W_\mu^a + i\frac{g'}{2}B_\mu\right)\Phi,$$
$$D_\mu \Delta = \partial_\mu \Delta + i\frac{g}{2}[\tau^a W_\mu^a, \Delta] + ig' B_\mu \Delta. \tag{4}$$

The fermion composition of the model is augmented by an extra generation of quarks and leptons ($SU(2)_L$ doublets and right-handed singlets):

$$Q_L^4 = \begin{pmatrix} t' \\ b' \end{pmatrix}_L, \quad t'_R, b'_R; \quad L_L^4 = \begin{pmatrix} \nu_{\tau'} \\ \tau' \end{pmatrix}_L, \quad \tau'_R, \nu'_R. \tag{5}$$

The Yukawa interaction for the Higgs fields is given by

$$\mathcal{L}_Y = -[\bar{Q}_L^i Y_d^{ij} \Phi d_R^j + \bar{Q}_L^i Y_u^{ij} \tilde{\Phi} u_R^j + \bar{L}_L^i Y_e^{ij} \Phi e_R^j + \bar{L}'_L Y_{\nu'} \tilde{\Phi} \nu'_R + \text{H.c.}] + h_{ij} \overline{L_L^{ic}} i\tau_2 \Delta L_L^j + \text{H.c.}, \tag{6}$$

where Φ, $\tilde{\Phi} = i\tau_2 \Phi^*$, $Y_{u,d,e,\nu'}$ are 4×4 complex matrices, and h_{ij} is a 4×4 complex symmetric Yukawa matrix. The triplet field Δ carries lepton number 2. The most general form of the Higgs potential involving the doublet Φ and triplet Δ under the gauge symmetry is given by [36]

$$V(\Phi, \Delta) = m_\Phi^2 \Phi^\dagger \Phi + M^2 \text{Tr}(\Delta^\dagger \Delta) + [\mu \Phi^T i\tau_2 \Delta^\dagger \Phi + \text{H.c.}] + \frac{\lambda}{4}(\Phi^\dagger \Phi)^2 + \lambda_1 (\Phi^\dagger \Phi)\text{Tr}(\Delta^\dagger \Delta) + \lambda_2 [\text{Tr}(\Delta^\dagger \Delta)]^2 + \lambda_3 \text{Tr}[(\Delta^\dagger \Delta)^2] + \lambda_4 \Phi^\dagger \Delta \Delta^\dagger \Phi, \tag{7}$$

where m_Φ and M are the mass-dimension real parameters, μ is the lepton-number violating parameter of mass dimension (which can be complex but is taken to be real here), and λ, λ_1-λ_4 are dimensionless real coupling constants.

From the stationary conditions at the vacuum for (v_Φ, v_Δ), we obtain

$$\frac{\partial V(\Phi, \Delta)}{\partial v_\Phi} = 0, \qquad \frac{\partial V(\Phi, \Delta)}{\partial v_\Delta} = 0,$$

yielding conditions for m_Φ^2, M^2:

$$m_\Phi^2 = \frac{1}{2}\left[-\frac{v_\Phi^2 \lambda}{2} - v_\Delta^2(\lambda_1 + \lambda_4) + 2\sqrt{2}\mu v_\Delta\right], \tag{8}$$

$$M^2 = M_\Delta^2 - \frac{1}{2}[2v_\Delta^2(\lambda_2 + \lambda_3) + v_\Phi^2(\lambda_1 + \lambda_4)], \quad \text{with}$$
$$M_\Delta^2 \equiv \frac{v_\Phi^2 \mu}{\sqrt{2} v_\Delta}, \tag{9}$$

which can be used to eliminate m_Φ^2 and M^2. The mass matrices for the scalar bosons can be diagonalized by rotating the scalar fields as

$$\begin{pmatrix} \varphi^\pm \\ \Delta^\pm \end{pmatrix} = \begin{pmatrix} \cos\beta_\pm & -\sin\beta_\pm \\ \sin\beta_\pm & \cos\beta_\pm \end{pmatrix} \begin{pmatrix} w^\pm \\ H^\pm \end{pmatrix},$$
$$\begin{pmatrix} \chi \\ \eta \end{pmatrix} = \begin{pmatrix} \cos\beta_0 & -\sin\beta_0 \\ \sin\beta_0 & \cos\beta_0 \end{pmatrix} \begin{pmatrix} z \\ A \end{pmatrix},$$
$$\begin{pmatrix} \varphi \\ \delta \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} h \\ H \end{pmatrix}, \tag{10}$$

where we defined the mixing angles as

$$\tan\beta_\pm = \frac{\sqrt{2} v_\Delta}{v_\Phi}, \qquad \tan\beta_0 = \frac{2v_\Delta}{v_\Phi},$$
$$\tan 2\alpha = \frac{4v_\Delta}{v_\Phi} \frac{v_\Phi^2(\lambda_1 + \lambda_4) - 2M_\Delta^2}{v_\Phi^2 \lambda - 2M_\Delta^2 - 4v_\Delta^2(\lambda_2 + \lambda_3)}. \tag{11}$$

In addition to the three Goldstone bosons $w^\pm$ and z, which give mass to the gauge bosons, there are seven physical mass eigenstates, $H^{\pm\pm}$, $H^\pm$, A, H and h. The masses of these physical states are expressed in terms of the parameters in the Lagrangian as

$$m_{H^{++}}^2 = M_\Delta^2 - v_\Delta^2 \lambda_3 - \frac{\lambda_4}{2} v_\Phi^2, \tag{12}$$

$$m_{H^+}^2 = \left(M_\Delta^2 - \frac{\lambda_4}{4} v_\Phi^2\right)\left(1 + \frac{2v_\Delta^2}{v_\Phi^2}\right), \tag{13}$$

$$m_A^2 = M_\Delta^2 \left(1 + \frac{4v_\Delta^2}{v_\Phi^2}\right), \tag{14}$$

$$m_H^2 = \mathcal{M}_{11}^2 \sin^2\alpha + \mathcal{M}_{22}^2 \cos^2\alpha - \mathcal{M}_{12}^2 \sin 2\alpha, \tag{15}$$

$$m_h^2 = \mathcal{M}_{11}^2 \cos^2\alpha + \mathcal{M}_{22}^2 \sin^2\alpha + \mathcal{M}_{12}^2 \sin 2\alpha, \quad (16)$$

where $\mathcal{M}_{11}^2$, $\mathcal{M}_{22}^2$ and $\mathcal{M}_{12}^2$ are the elements of the mass matrix $\mathcal{M}_{ij}^2$ for the CP-even scalar states in the (φ, δ) basis, which are given by

$$\mathcal{M}_{11}^2 = \frac{v_\Phi^2 \lambda}{2}, \quad (17)$$

$$\mathcal{M}_{22}^2 = M_\Delta^2 + 2v_\Delta^2(\lambda_2 + \lambda_3), \quad (18)$$

$$\mathcal{M}_{12}^2 = -\frac{2v_\Delta}{v_\Phi} M_\Delta^2 + v_\Phi v_\Delta(\lambda_1 + \lambda_4). \quad (19)$$

The masses of the W and Z bosons are obtained at the tree level as

$$m_W^2 = \frac{g^2}{4}(v_\Phi^2 + 2v_\Delta^2),$$
$$m_Z^2 = \frac{g^2}{4\cos^2\theta_W}(v_\Phi^2 + 4v_\Delta^2). \quad (20)$$

The electroweak ρ parameter is defined at the tree level as

$$\rho \equiv \frac{m_W^2}{m_Z^2 \cos^2\theta_W} = \frac{1 + \frac{2v_\Delta^2}{v_\Phi^2}}{1 + \frac{4v_\Delta^2}{v_\Phi^2}}. \quad (21)$$

As the experimental value of the ρ parameter is near unity, v_Δ^2/v_Φ^2 is required to be much smaller than unity at the tree level. Note, in fact, that in the HTM4 the ρ parameter is less than 1, which may contradict the particle data group fit [13] $\rho_0 = 1.0008^{+0.0017}_{-0.0007}$ obtained from a global fit including the direct search limits on the standard Higgs boson. However, at the 2σ level, $\rho_0 = 1.0004^{+0.0029}_{-0.0011}$ [13], which is compatible with $\delta\rho < 0$. Relaxing the direct limit on the Higgs mass yields $\rho_0 = 1.0008^{+0.0017}_{-0.0010}$, again compatible with $\delta\rho < 0$, which implies an upper bound on v_Δ of the order 2.5–4.6 GeV. Thus, a $v_\Delta \sim \mathcal{O}(1)$ GeV would safely fit the constraints. Barring accidental cancellations, the mixing angles are small, the state h behaves mostly as the SM Higgs boson, while the other states are almost entirely components of the triplet field. Note that the smallness of v_Δ/v_Φ ensures that the mixing angles $\beta_\pm$ and β_0 are close to 0, but given the complex expression defining it, it is worthwhile to note that α, the mixing angle between the doublet and triplet neutral Higgs bosons, remains undetermined [37].

In the HTM4, tiny Majorana neutrino masses are generated by the Yukawa interaction with the VEV of the triplet field, which is proportional to the lepton number violating coupling constant μ as

$$(m_\nu)_{ij} = \sqrt{2} h_{ij} v_\Delta = h_{ij} \frac{\mu v_\Phi^2}{M_\Delta^2}. \quad (22)$$

If $\mu \ll M_\Delta$ the smallness of the neutrino masses are explained by the type-II seesaw mechanism.

III. OBLIQUE CORRECTIONS

The contributions to the oblique parameters can in general be written as

$$S = \frac{4c_W^2 s_W^2}{\alpha_{em}}\left[\frac{\Pi_{\gamma\gamma}^{1PI}(m_Z^2) - \Pi_{\gamma\gamma}^{1PI}(0)}{m_Z^2} + \frac{c_W^2 - s_W^2}{c_W s_W}\frac{\Pi_{Z\gamma}^{1PI}(m_Z^2) - \Pi_{Z\gamma}^{1PI}(0)}{m_Z^2} - \frac{\Pi_{ZZ}^{1PI}(m_Z^2) - \Pi_{ZZ}^{1PI}(0)}{m_Z^2}\right],$$
$$T = \frac{1}{\alpha_{em}}\left[\frac{\Pi_{ZZ}^{1PI}(0)}{m_Z^2} - \frac{\Pi_{WW}^{1PI}(0)}{m_W^2} + \frac{2s_W}{c_W}\frac{\Pi_{Z\gamma}^{1PI}(0)}{m_Z^2} + \frac{s_W^2}{c_W^2}\frac{\Pi_{\gamma\gamma}^{1PI}(0)}{m_Z^2}\right],$$
$$U = \frac{4s_W^2}{\alpha_{em}}\left[s_W^2\frac{\Pi_{\gamma\gamma}^{1PI}(m_Z^2) - \Pi_{\gamma\gamma}^{1PI}(0)}{m_Z^2} + 2c_W s_W\frac{\Pi_{Z\gamma}^{1PI}(m_Z^2) - \Pi_{Z\gamma}^{1PI}(0)}{m_Z^2} + c_W^2\frac{\Pi_{ZZ}^{1PI}(m_Z^2) - \Pi_{ZZ}^{1PI}(0)}{m_Z^2} - \frac{\Pi_{WW}^{1PI}(m_W^2) - \Pi_{WW}^{1PI}(0)}{m_W^2}\right]. \quad (23)$$

The expressions for $\Pi_{ZZ}^{1PI}(p^2), \Pi_{Z\gamma}^{1PI}(p^2), \Pi_{WW}^{1PI}(p^2)$ and $\Pi_{\gamma\gamma}^{1PI}(p^2)$ for the HTM are given in terms of Passarino-Veltman functions in Appendix B of Ref. [38], and the contributions to the oblique corrections due to the sequential fourth generation can be found in Ref. [39]. Note that both the new scalar sector (with nontrivial mixing amongst themselves and its modified couplings to the weak gauge bosons) and the fourth generation fermions will lead to significant and nonvanishing contributions to the oblique electroweak corrections. It is well known that in the limit of degenerate isospin multiplets (in this case the fourth generation family of leptons and quarks) there is a positive contribution $\Delta S = 0.21$, which can be significantly altered once the HTM contributions are included. In fact, as we show later through our analysis, the HTM contributions cancel the large positive contributions coming from the fourth generation even with the degeneracy not lifted between the isodoublets. We also find that a large mass

splitting in the fermion isodoublets which would otherwise have not been preferred because of the large ΔT corrections can still be allowed once the HTM and SM4 contributions are combined.

IV. ANALYSIS FRAMEWORK

The main motivation in this analysis is to salvage the fourth generation chiral fermions. The HTM4 model contains, in addition to the SM particle content, a complex triplet scalar along with a sequential fourth generation of chiral fermions. The free parameters in the theory are the fourth generation fermion masses [we neglect the Cabibbo-Kobayashi-Maskawa (CKM) mixings], the independent coefficients in the potential of the Higgs sector [see Eq. (7))] and the VEVs for the neutral components of the scalar multiplets. To limit the available number of parameters we have used the various constraints on the coefficients of the Higgs potential (Sec. B). This reduces the effective number of free parameters. In addition, we consider both the masses of τ' and $\nu_{\tau'}$ to be fixed at 250 GeV. The reason is the following: the strongest constraint on the fourth generation lepton masses comes from the Z-boson width, forcing the fourth neutrino to be heavier than $M_Z/2$. The most stringent bounds on the fourth generation lepton masses come from L3 at the LEP [40], obtained from e^+e^- collisions at $\sqrt{s} = 200$ GeV:

$$m_{l'} > 110.8\text{GeV}, \qquad m_{\nu'} > \begin{cases} 90.3\ \text{GeV} & \text{for Dirac}\,\nu \\ 80.5\ \text{GeV} & \text{for Majorana}\,\nu, \end{cases} \tag{24}$$

rendering our choice of parameters quite conservative. We also note that large electroweak corrections for the $h \to gg$ decay mode can be moderated to small values with the above choice (or lower) of the fourth generation lepton masses [41].

The method of our analysis is roughly based on the following steps. Using HIGLU [42], we calculate the next-to-next-to-leading order cross sections for Higgs production in the gluon fusion channel, obtained by varying the fourth generation quark masses between 550 GeV and 750 GeV. Our choice of masses is based on the results of searches by CMS and ATLAS. But we allow for a softening of the limits, as these are obtained using SM4 specific assumptions. Direct searches at ATLAS and CMS always assume that a single specific decay has 100% branching ratio, usually $t' \to bW$ and $b' \to tW$. While constraints on the mixing between the first and second generations with the fourth generation can be extracted from flavor-changing neutral current bounds in $K^0 - \bar{K}^0$ and $D^0 - \bar{D}^0$, the size of the CKM4 matrix elements is still allowed to be larger than the smallest matrix elements in the usual CKM matrix [43], $|V_{ub'}|, |V_{t'd}|, |V_{cb'}| \lesssim 0.04$. Additionally, the mixing between the third and fourth generations is very weakly constrained, as there is a weak limit from the single top production, $|V_{tb}| > 0.89 \pm 0.07$ [13], and this bound is stronger than that obtained using the electroweak precision data. Though, while it seems likely that the fourth generation quarks will decay into the third generation, this need not be so, and not with 100% branchings, which would modify the limits on the masses. The fourth generation fermions can also decay to a Higgs boson via flavor violating couplings to Higgs, which can be introduced by including certain higher dimensional operators [44]. This can further lead to weakening of the constraints on the fourth generation quark masses. These considerations allow us to vary the mass of fourth generation quarks from values much below the current experimental bounds, which assume very specific decay patterns for the heavy quarks in their analyses.

We find the (gluon fusion) production cross section to be ~195 pb (~153 pb) at 8 TeV (7 TeV), which is about a factor of 10 larger than its SM expectation. This enhancement is known, and is one of the reasons why the fourth generation chiral fermions is ruled out without additional new physics effects. Next, we calculate the decay widths of the Higgs boson in its various decay modes using HDECAY [45]. Here also, we vary the masses of the fourth generation quark masses in the 550–750 GeV range. We include the electroweak corrections (EW) coming from the fourth generation contributions in all the tree level decay processes as well as the loop induced gluon mode. However, we have chosen to ignore the electroweak corrections to $h \to \gamma\gamma, Z\gamma$ decay widths in our analysis, as there are additional particles ($H^\pm$ and $H^{\pm\pm}$) contributing in the loop and therefore the EW contributions will be severely altered from what is computed in the literature. While the electroweak radiative corrections to the $gg \to h$ process are significant [21], for a specific choice of fermion masses, $m_{b'} = m_{t'} + 50\,\text{GeV} = m_{\ell'} = m_{\nu'} \sim 600\,\text{GeV}$, they lead to an increase (decrease) of the cross section at low (high) Higgs masses, $M_h \approx 120(600)$ GeV, by ≈12% [43]. We also note that the EW corrections in the $gg \to h$ mode can be kept within 5% with the appropriate choice for the mass of the fourth generation leptons [41]. However, for the decays $h \to f\bar{f}$ and $h \to VV$, the $\mathcal{O}(G_F m_{f'}^2)$ terms, implemented by multiplying the couplings g_{hXX} by the electroweak correction term $1 + \delta_{ew}^X$, have been included in HDECAY [45,46]. Results from the precise calculations indicate that the approximation of including only the leading terms also works very well in this case [21].

We then separately scan the parameter space including the Higgs and the electroweak sector. We impose the existence of a boson of mass ~125 GeV, consistent with the particle discovered at the LHC, and which has properties similar to the SM Higgs (whether it is the SM Higgs or not will only be revealed by knowing its properties to a more accurate extent). In our model, which has two CP-even neutral scalar Higgs bosons, we constrain the mass of one of these to be in the range 123–127 GeV. We keep this small window in the mass in

order to account for the uncertainties in the mass measurements in the various channels of the Higgs decay. We constrain the second CP-even neutral Higgs to be in the range 97–99 GeV. Our inspiration in choosing a Higgs boson of such a mass is the 2.3σ excess seen at 98 GeV by the LEP [47,48], which has still not been ruled out experimentally and is outside LHC and Tevatron sensitivity. From the parameters v_Δ, μ, λ, λ_1, λ_2, λ_3 and λ_4, we can calculate the masses of the two neutral CP-even Higgs bosons, m_h and m_H, the CP-odd Higgs, m_A, and of the four charged Higgs bosons, $m_{H^\pm}$ and $m_{H^{\pm\pm}}$. From the above parameters, we also obtain the four mixing angles, β_0, $\beta_\pm$ and α. (See Sec. A where we give the formulas relating the above seven parameters to the Higgs masses and the four mixing angles). To reduce the enhanced cross section of the Higgs (h) production in the gluon fusion channel, we let $\cos^2\alpha$ vary between $1/12$ and $1/2$.[3] In addition, we let both μ and λ vary, such that the constraints on their magnitudes (as given in Sec. B) are satisfied. We allow λ_1 to vary around a small positive number, but do not impose an upper bound. We apply a simplifying assumption by choosing $\lambda_2 = \lambda_3$. Also, we choose negative values of λ_4, yielding the mass hierarchy $m_{H^{\pm\pm}} > m_{H^\pm}$. (Should we have chosen λ_4 to be positive, we would have obtained a different hierarchy in the masses of the charged Higgs bosons).

To check with the current Higgs data and see what portion of the parameter space is allowed, we construct the theoretical signal strength in the ith channel as defined below:

$$\mu^i = \frac{R_i^{\rm prod} \times R_i^{\rm decay}}{R^{\rm width}}, \tag{25}$$

where μ^i is the theoretically computed signal strength and $R_i^{\rm prod}$, $R_i^{\rm decay}$ and $R^{\rm width}$ are the factors modifying the production cross section, decay width in the ith channel and the total decay width, respectively. The signal strengths provided by the experimental collaborations are given by $\hat{\mu}^i = \sigma_i^{\rm obs}/\sigma_i^{\rm SM}$, with their 1σ uncertainties given by σ_i. Here, $\sigma_i^{\rm obs}$ denotes the observed signal cross section for a particular Higgs mass, whereas $\sigma_i^{\rm SM}$ is the signal cross section for an SM Higgs boson of the same mass. In order to obtain μ^i for each decay channel, we compute the modifications in the production cross sections from the various modes, the modifications in the partial decay widths in the various channels, as well as the change in the total decay width. The modifications in the production cross sections are as follows:

(i) The modification in the gluon fusion production channel, R_{GF}, is a function of the masses of the fourth generation chiral fermions and the mixing angle α.
(ii) The factors modifying the production cross sections in the weak boson fusion channel or in the associated production mode, viz. R_{VBF}, R_{Wh} and R_{Zh}, are functions of v_ϕ, v_Δ and α.
(iii) The modification in the $t\bar{t}h$ production mode, $R_{t\bar{t}h} = \cos^2\alpha$.

Similarly, the factors modifying the decay widths are:

(i) $$R_{h\to\gamma\gamma} = \frac{|\sum_{f=t,b,\tau,t',b',\tau'} N_c^f Q_f^2 g_{hff} A_{1/2}^h(\tau_f) + g_{hWW} A_1^h(\tau_W) + g_{hH^\pm H^\mp} A_0^h(\tau_{H^\pm}) + 4 g_{hH^{\pm\pm}H^{\mp\mp}} A_0^h(\tau_{H^{\pm\pm}})|^2}{|\sum_{f=t,b,\tau} N_c^f Q_f^2 g_{hff} A_{1/2}^h(\tau_f) + g_{hWW} A_1^h(\tau_W)|^2}$$

Please see Sec. A1 for more details on the formulas.

(ii) $R_{h\to ZZ^*} = (v_\varphi \cos\alpha + 4 v_\Delta \sin\alpha)^2/v^2$.
(iii) $R_{h\to WW^*} = (v_\varphi \cos\alpha + 2 v_\Delta \sin\alpha)^2/v^2$.
(iv) After computing the fermionic decay widths using HDECAY by varying the fourth generation quark masses, we further modify them by a factor $(\cos\alpha/\cos\beta_\pm)^2$. We then divide these by the corresponding SM decay widths to obtain $R_{h\to b\bar{b}}$ and $R_{h\to\tau\bar{\tau}}$.
(v) We also compute $R_{h\to gg}$, $R_{h\to\mu^+\mu^-}$, $R_{h\to c\bar{c}}$, $R_{h\to s\bar{s}}$ and $R_{h\to Z\gamma}$ in order to compute $R^{\rm width}$, i.e. the modification in the total decay width of h.

After constructing the μ^is for the full parameter space, we compare these with the 2σ allowed ranges of $\hat{\mu}^i$ as given by the experimental collaborations (see Table I).[4] This gives us a reduced parameter space. We then proceed to check if this reduced parameter space satisfies the S, T and U bounds within the 1σ uncertainty limits [50]. For consistency, we have checked our results for oblique parameters against those in the SM [51]. We performed our analysis for two benchmark values of v_Δ ($v_\Delta = 1,\ 3$ GeV). We found that, after imposing all conditions, we could not find any surviving parameter points for $v_\Delta = 1$ GeV within the 1σ range for ΔS and ΔT combined, although a significant region is allowed

[3]Naively, one might expect that a $\cos^2\alpha \sim 1/10$ will compensate for the $\mathcal{O}(10)$ enhancement in the gluon fusion initiated cross section. But, as is evident from some of the decay widths, the α enters in a nontrivial manner and not just as an overall $\cos^2\alpha$ factor. Hence, we keep this range instead of a single fixed value because our aim is to ensure that the value of $\sigma\times$ branching ratio (BR) is close to its SM counterpart.

[4]We deliberately avoid an involved statistical analysis (such as χ^2) because of the large number of parameters and limited number of data points as we have only considered the inclusive signal strengths (apart from the $b\bar{b}$ channel where the results correspond to the associated production mode only) for the various channels given by ATLAS and CMS [49].

TABLE I. Data set used in our analysis, with the values of $\hat{\mu}_i$ in various channels and their 1σ uncertainties as reported by the ATLAS [52,53] and CMS collaborations [54–57].

Channel	$\hat{\mu}$	Experiment	Energy in TeV (Luminosity in fb^{-1})
$h \to \gamma\gamma$	$1.55^{+0.33}_{-0.28}$	ATLAS	$7(4.8) + 8(20.7)$
$h \to \gamma\gamma$	$0.78^{+0.28}_{-0.26}$	CMS	$7(5.1) + 8(19.6)$
$h \overset{ZZ^*}{\to} 4l$	$1.43^{+0.40}_{-0.35}$	ATLAS	$7(4.6) + 8(20.7)$
$h \overset{ZZ^*}{\to} 4l$	$0.93^{+0.29}_{-0.25}$	CMS	$7(5.1) + 8(19.7)$
$h \overset{WW^*}{\to} 2l2\nu$	$0.99^{+0.31}_{-0.28}$	ATLAS	$7(4.6) + 8(20.7)$
$h \overset{WW^*}{\to} 2l2\nu$	$0.72^{+0.20}_{-0.18}$	CMS	$7(4.9) + 8(19.4)$
$h \to b\bar{b}$	$0.20^{+0.70}_{-0.60}$	ATLAS (VH)	$7(4.7) + 8(20.3)$
$h \to b\bar{b}$	$1.00^{+0.50}_{-0.50}$	CMS (VH)	$7(5.1) + 8(18.9)$
$h \to \tau\bar{\tau}$	$1.4^{+0.50}_{-0.40}$	ATLAS	$8(20.3)$
$h \to \tau\bar{\tau}$	$0.78^{+0.27}_{-0.27}$	CMS	$7(4.9) + 8(19.7)$
$h \overset{WW^*}{\to} 2l2\nu$	$1.4^{+0.70}_{-0.60}$	ATLAS (VBF)	$7(4.6) + 8(20.7)$
$h \overset{WW^*}{\to} 2l2\nu$	$0.62^{+0.58}_{-0.47}$	CMS (VBF)	$7(4.9) + 8(19.4)$

by the Higgs data. Thus, the model is restricted to larger values of v_Δ. However, we find that at 3σ allowed ranges of ΔS and ΔT, lower values of $v_\Delta \simeq 1$ GeV could still give us an allowed region in the parameter space. Thus, we may claim that our choice for the allowed range in the S, T and U planes for oblique corrections is slightly more restrictive.

We must add that in our analysis, we took the same cut-efficiencies for all the production channels while computing the signal strengths. This should not affect our results as the cross section coming from the gluon fusion production mode dominates significantly over the other production modes. The following formulas were used in combining the theoretically computed μ^i values for the $Wh \to l\nu b\bar{b}$, $Zh \to l^+l^-b\bar{b}$ and $Zh \to \nu\bar{\nu}b\bar{b}$:

$$\frac{1}{\bar{\sigma}^2} = \sum_i \frac{1}{\sigma_i^2},$$
$$\frac{\bar{\mu}}{\bar{\sigma}^2} = \sum_i \frac{\mu_i}{\sigma_i^2}. \tag{26}$$

These yield the combined 1σ uncertainties and the combined signal strengths. Since the experimental collaborations have reported a single signal strength for the $h \to b\bar{b}$ channel in the associated production mode, such combinations have to be included.

V. RESULTS

We now present our results for the parameter scan of the HTM4 which is consistent with the Higgs data and the oblique correction constraints. We find that there is a significant region in the parameter space which ensures that the otherwise constrained model for fourth generation chiral fermions is still allowed if we consider the effects of a triplet Higgs sector, provided it has a nontrivial mixing with the scalar doublet in the SM. We have illustrated our results through some scatter plots for the various coefficients of the scalar potential giving rise to the aforementioned nontrivial mixing in the Higgs sector and gives us a ∼125 GeV scalar consistent with the LHC data. In Fig. 1, we show the allowed parameter space in the $\lambda_4 - \lambda$, $\lambda_4 - \mu$ and $\mu - \lambda$ planes for $m_{t'}, m_{b'}$ lying in the region 550–600 GeV. Note that although our scan over the fourth generation quark masses is between 550–750 GeV, the Higgs data seems to more severely constrain the fourth generation quark mass above 600 GeV, at least within the 2σ uncertainty limits of all respective $\hat{\mu}^i$. This is found to happen because the EW corrections to the $h \to ZZ^*, WW^*$ reduce the branching fractions of these modes by a large amount for higher values of the quark masses, which therefore affects the Higgs data significantly. We therefore allow for a 3σ deviation in the experimentally observed signal strengths and find that we do get a viable region in the parameter space for the higher values of the fourth generation quark masses.

We must note that, in all of these plots, the other parameters are also varied. This way, we have shown the projection of the allowed parameter space in the above three planes. We find that λ has an allowed range of 0.35–0.37 and μ has an approximate allowed range of 0.96–1.03, when we consider the 2σ Higgs data constraints (as explained in the previous section). On the other hand, given our choice of negative λ_4, the allowed region from the Higgs data varies between −0.6 and −0.2. But as we can see from each of these plots, when we calculate the oblique corrections in the HTM4 model and impose the S, T and U constraints at 1.25σ[5]deviations from the respective central values, the allowed region reduces considerably. It is worth noting that we have been conservative in allowing only 1.25σ deviations in ΔS and ΔT values, and that we are still able to achieve a viable parameter space with larger mass splittings in the fourth generation quark masses allowed by oblique corrections when compared to the conventional SM4. In fact, we find that for nondegenerate masses for $\nu_{\tau'}$ and τ' (allowing them to have a certain mass splitting) it is possible to obtain a significantly larger parameter space still allowed by the oblique parameter constraints. So, apart from varying the fourth generation quark masses, we could in principle have varied the fourth generation lepton masses also (while keeping the EW corrections small) and obtained a larger parameter space.

If we increase the masses of the fourth generation quarks by allowing them to vary in the region 550–750 GeV, the region of the parameter space that survives oblique

[5]We find that for 1σ constraints in S, T and U, we do not get any viable parameter points which also satisfy the Higgs data within 2σ standard deviations.

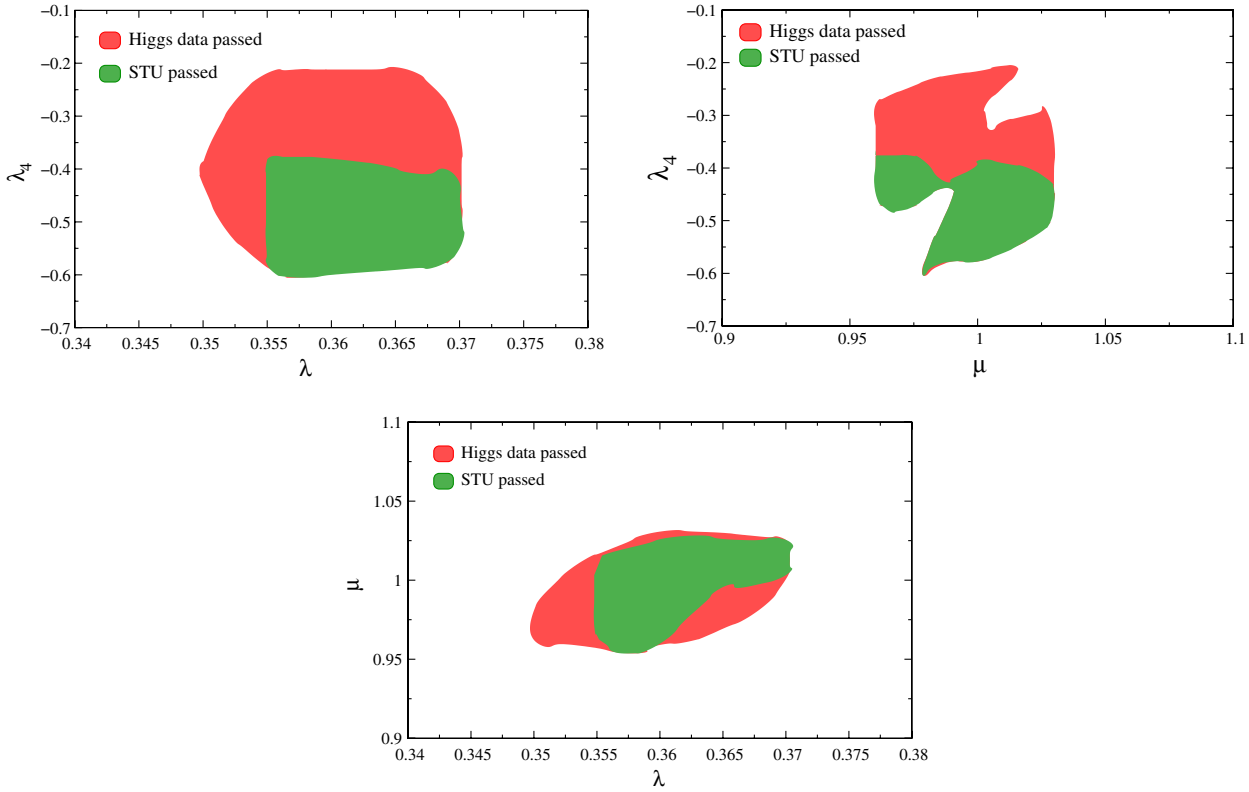

FIG. 1 (color online). Scatter plots for the allowed regions in the parameter space: λ_4 as a function of λ (top left); λ_4 as a function of μ (top right); μ as a function of λ (bottom). The full shaded (red and green) regions are allowed by the Higgs data while only the green region is allowed by the S, T and U constraints at 1.25σ. In all the plots, apart from the axes shown, all the other parameters were also varied over their allowed ranges. Note that we varied $m_{t'}, m_{b'} \simeq 550$–$600$ GeV while higher masses were disallowed from the Higgs data within 2σ standard deviations.

corrections (1σ) and the Higgs data (3σ) changes. In Fig. 2, we revisit the allowed parameter space in the $\lambda_4 - \lambda$, $\lambda_4 - \mu$ and $\mu - \lambda$ planes for heavier fourth generation masses. As before, in all of these plots the other parameters are also varied. Thus, the projection of the allowed parameter space in the above three planes shows that λ has an allowed range extending beyond 0.37 and μ has an approximate allowed range of 0.94–1.06, when we consider the Higgs data constraints to lie within 3σ errors. We also find viable regions of the parameter space with S, T and U constraints now at only 1σ.

As the survival of the fourth generation quarks in the current framework is closely related to the nontrivial mixing in the neutral scalar sector of our model, it becomes imperative to consider the allowed values of the mixing angle (α) which also satisfies the various constraints. In Fig. 3, we show the allowed parameter space in the $\sin\alpha$-λ and $\sin\alpha$-μ plane. Even though α is a function of λ and μ [see Eq. (11)], these plots help us to emphasize the importance of α. As discussed before, we not only require α to suppress the hugely enhanced cross section of the Higgs production from gluon fusion, but α also determines the mass spectrum for the neutral scalars. Note that the scalar mass spectrum consistent with a ~125 GeV Higgs boson is allowed over a range of $\sin\alpha$ lying between ~0.7–0.95. We find this a good place to discuss the features of the scalar mass spectrum that we obtain in our scan over the parameter space for $v_\Delta = 3$ GeV. Note that we have already specified the requirement for the mass of the two CP-even neutral scalars to be $m_H \sim 97$–99 GeV and $m_h \sim 123$–127 GeV (Sec. IV). The other scalars in the model are the pseudoscalar, singly and doubly charged scalars. For most of the allowed parameter values, the pseudoscalar has mass in the range $116 \leq m_A \leq 123$ GeV as it has a linear dependence on the choice of μ once the VEVs are fixed. The choice of λ_4 already decides the mass hierarchy for the charged scalars. The singly charged scalars lie within the mass range of ~118–214 GeV while the doubly charged scalars are in the mass range of ~125–276 GeV when satisfying the Higgs mass constraints. We find that $\sin\alpha$ ranges roughly between 0.85 and 0.93 (red region) solely from the Higgs data. This means that $\cos^2\alpha$ varies between 0.14 and 0.28. The allowed spectrum shrinks considerably once the oblique parameter corrections are included, yielding the modified ranges for the charged scalars as $m_{H^\pm} \sim 152$–170 GeV and $m_{H^{\pm\pm}} \sim 178$–208 GeV GeV. This range features in the plots for mass splittings in Fig. 4 and with slight variations

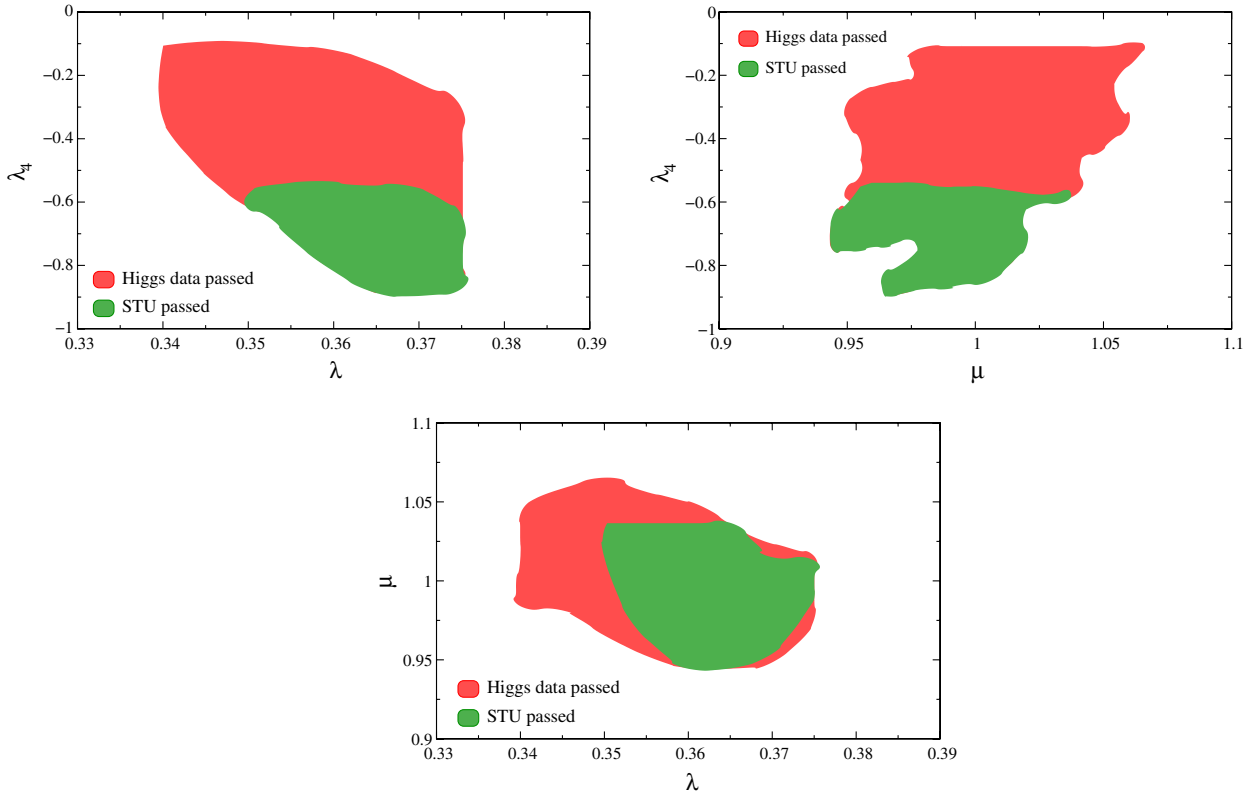

FIG. 2 (color online). Scatter plots for the allowed regions in the parameter space : λ_4 as a function of λ (top left); λ_4 as a function of μ (top right); μ as a function of λ (bottom). The full shaded (red and green) regions are allowed by the Higgs data while only the green region is allowed by the S, T and U constraints at 1σ. In all the plots, apart from the axes shown, all the other parameters were also varied over their allowed ranges. Note that here we have varied $m_{t'}, m_{b'} \simeq 550$–$750$ GeV while the allowed region from the Higgs data was increased to lie within 3σ errors.

(because of the weaker constraints being used from Higgs data) in Fig. 5. The other notable parameters that are worth mentioning are λ_1, which prefers a large value to allow for large mixing in the neutral scalars, while values of $0 < \lambda_{2,3} < 5$ are preferred for the final allowed regions illustrated in our plots. It is worth pointing out that the triplet multiplet will require a very weak Yukawa coupling to the lepton doublets (since v_Δ is large) to generate the tiny neutrino masses [see Eq. (22)], and therefore the doubly charged scalars in the model will not have the usual dominant leptonic decay modes which have been extensively used by experimentalists to put limits on their mass, while the singly charged Higgs bosons masses are within their experimental limits [13]. This ensures that the mass spectrum we obtain is safe from existing collider limits on such scalars. A detailed analysis of the collider signals for such scalars is left for future work.

Both the fourth generation fermions and the Higgs bosons in this model (HTM4) significantly affect the oblique corrections through their contributions to the self energies of the gauge bosons. We highlight the mass splittings for the fourth generation quarks and the charged scalars in the model which give a parameter space that respects the Higgs data limits as well as the constraints from the oblique corrections. Note that for our analysis, we have assumed degenerate fourth generation leptons ($m_{\tau'} = m_{\nu'} = 250$ GeV). We illustrate our findings through the contour plots of ΔS and ΔT in the plane of $m_{t'} - m_{b'}$ versus $m_{H^{\pm\pm}} - m_{H^\pm}$ (see Figs. 4 and 5). The colour coded index in each case shows the values of ΔS (for the left curve) and ΔT (for the right curve). The results are plotted for $m_{t'}, m_{b'} \sim 550$–600 GeV in Fig. 4, and for $m_{t'}, m_{b'} \sim 550$–750 GeV in Fig. 5. While in Figs. 1 and 2 the parameter space varied only slightly when increasing the fourth generation masses, the change is much more noticeable in here, where a larger region of the parameter space survives for $m_{t'}, m_{b'} \sim 600$–750 GeV. An interesting feature to note is that even though the contributions in HTM4 for ΔS are found to be only positive, the contributions of the model in ΔT are found to be both positive and negative. The negative contributions to the ΔT come from the bosonic loops and from the extended scalar sector, with the nontrivial mixing also playing a significant role. From the plots, we find that $m_{t'}$ can be either greater than or less than $m_{b'}$ for the allowed parameter space. On the other hand, $m_{H^{\pm\pm}}$ is always greater than the $m_{H^\pm}$, due to our original choice requiring λ_4 to be negative. Clearly we find that a larger mass splitting is allowed in HTM4 when

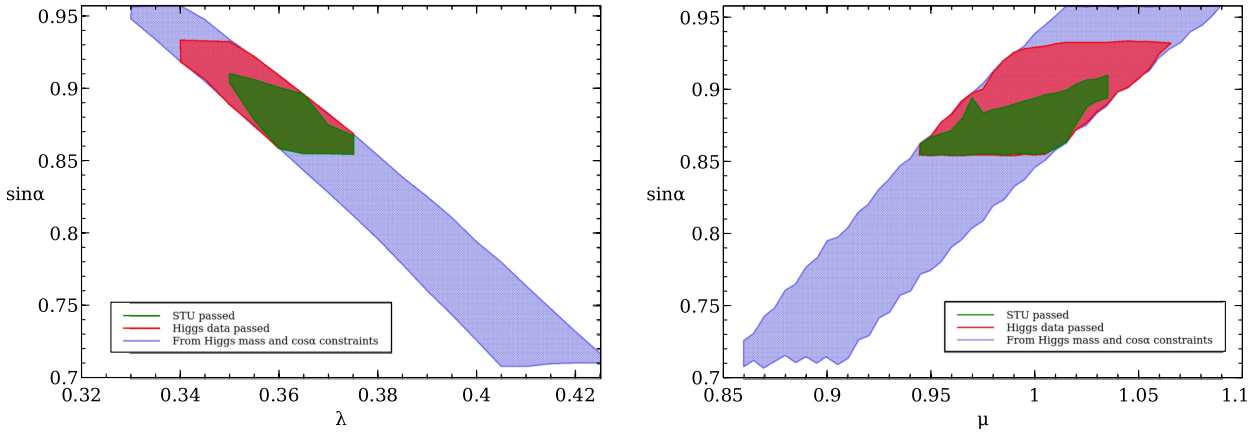

FIG. 3 (color online). Scatter plots for the allowed regions in the parameter space: $\sin\alpha$ as a function of λ (left); $\sin\alpha$ as a function of μ (right). The blue shaded regions are allowed by the Higgs mass constraints, the red shaded regions are then allowed by the Higgs data (3σ), while the green shaded regions are then allowed by the S,T and U constraints (1σ). In both the plots, apart from the axes shown, all the other parameters were also varied over their allowed ranges.

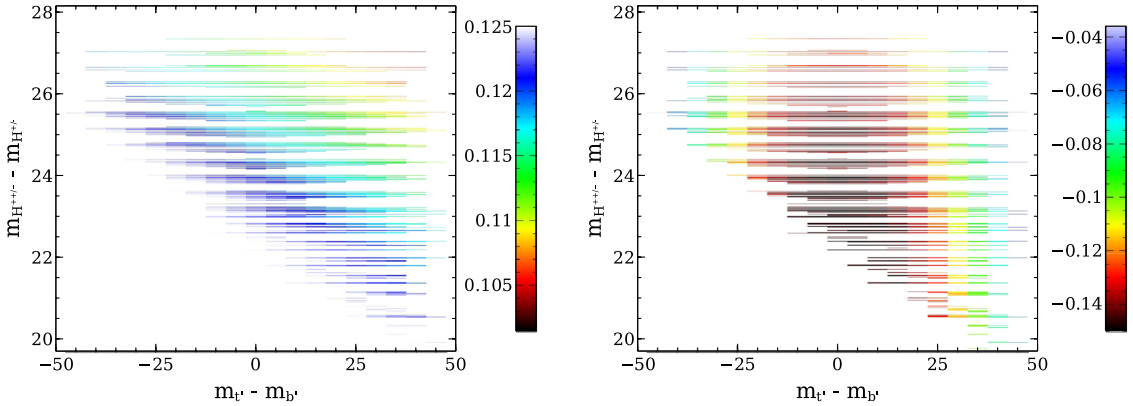

FIG. 4 (color online). Contour plots for ΔS (left) and ΔT (right) with $m_{t'} - m_{b'}$ and $m_{H^{\pm\pm}} - m_{H^{\pm}}$ varied. We vary $m_{t'}, m_{b'} \sim 550$–600 GeV, which is consistent with Higgs data within errors of 2σ. The color index at the right side of each plot shows the variation of ΔS (left) and ΔT (right).

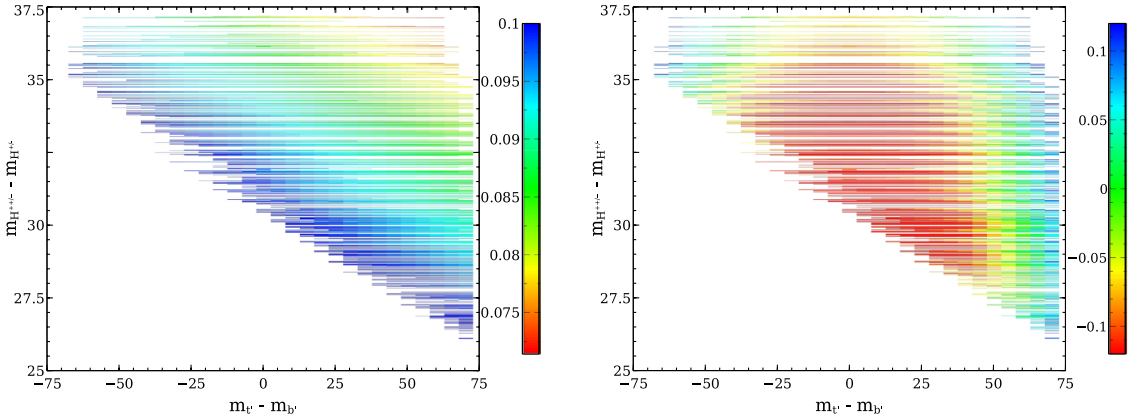

FIG. 5 (color online). Contour plots for ΔS (left) and ΔT (right) with $m_{t'} - m_{b'}$ and $m_{H^{\pm\pm}} - m_{H^{\pm}}$ varied. We vary $m_{t'}, m_{b'} \sim 550$–750 GeV, which is consistent with Higgs data within errors of 3σ. The color index at the right side of each plot shows the variation of ΔS (left) and ΔT (right).

compared to the SM4 which is disfavored from Higgs data alone.

Note that the mixing angle α will suppress the production cross section of the ∼125 GeV Higgs boson in the VBF and VH channels. Although these channels, when considered inclusively along with the gluon fusion channel, are subdominant, we check the rates for them against the values shown in Table I for a representative point which has passed all the constraints discussed earlier. For the point $\sin\alpha \approx 0.87$, $\lambda = 0.365$, $\mu = 1.0$, $\lambda_4 = -0.79$, $m_{t'} = 600$ GeV, $m_{b'} = 650$ GeV in the parameter space, we find that $\frac{[\sigma\times \mathrm{BR}]^{\mathrm{SM4+triplet}}_{\mathrm{VBF}}}{[\sigma\times \mathrm{BR}]^{\mathrm{SM}}_{\mathrm{VBF}}} \simeq 0.062$ in the WW^* channel (2.11σ away from the combined central value of ATLAS and CMS) and $\frac{[\sigma\times \mathrm{BR}]^{\mathrm{SM4+triplet}}_{\mathrm{VH}}}{[\sigma\times \mathrm{BR}]^{\mathrm{SM}}_{\mathrm{VH}}} \simeq 0.243$ in the $b\bar{b}$ channel (1.16σ away from the combined central value of ATLAS and CMS). In fact, we find that for all the STU passed points, the VBF signal strengths for the WW^* channel are within $\sim 2.15\sigma$ while the VH signal strengths for the $b\bar{b}$ channel are within $\sim 1.35\sigma$ from their respective experimental signal strengths.

VI. SUMMARY AND CONCLUSION

An extension of the SM by a fourth sequential generation of quarks and leptons is one of the simplest beyond the SM scenarios. The fourth generation rectifies some of the shortcomings of the SM, and, in particular, it provides a dark matter candidate, ν'. Electroweak precision measurements (oblique and nonoblique corrections) favor small mass splittings between the fourth generation quarks, $m_{t'} - m_{b'} < M_W$, and flavor violating decays are restricted by the structure of the 3×3 CKM matrix.

However, for all its nice features, the fourth generation has all but been abandoned. This is mainly due to two experimental inputs. First, discovery of the Higgs boson and recent measurements of the Higgs production cross section and decay rates disfavor the SM with four generations. The Higgs production through gluon fusion is enhanced by the presence of new, $SU(3)_c$ interacting heavy particles, rendering it inconsistent with experimental results associated with the Higgs boson at ∼125 GeV. Second, direct limits on the masses of fourth generation quarks, all assuming that the decays $t' \to Wb$ and $b' \to Wt$ have 100% branching ratios, have pushed limits on the fourth generation quark masses into the 700–800 GeV region. This has serious implications on Yukawa couplings, putting perturbativity of the couplings in serious jeopardy.

In this work we show that relaxing the single Higgs boson requirement and slightly relaxing the experimental bounds (that is, allowing for alternative decays of the fourth generation fermions) saves the fate of the fourth generation. We choose to work in the Higgs triplet model, where the neutral triplet Higgs state is allowed to mix nontrivially with the SM doublet Higgs state. Unlike in two Higgs doublet models, in this model the Yukawa couplings are not proportional to the ratio of the doublet VEVs, lifting some of the pressure on the perturbativity of couplings. We allow for a second neutral Higgs boson, just below the sensitivity of both Tevatron and LHC, but consistent with the 2σ bump at the LEP in the 98–99 GeV region. We impose precision electroweak constraints on the model, constraints on the masses of the singly and doubly charged Higgs masses, as well as the requirement that the production and decay rates we obtain reproduce the data from ATLAS and CMS, within their errors. We scan the parameter space of the HTM4 model under these restricted conditions and show that some regions of the parameter space, *albeit* restricted but still significant, survive. Couplings in the Higgs potential are limited to small regions: λ has an allowed range between 0.34 and 0.375, μ has an approximate allowed range between 0.94 and 1.07, and λ_4, chosen to be negative, has an allowed range between −0.9 and −0.1. The precise value for $\sin\alpha$, which ranges roughly between 0.85 and 0.93 from the Higgs data at 3σ, is consistent with the other constraints. We find that $m_{t'}$ can be either larger or smaller than $m_{b'}$, and the parameter space for $m_{b'}, m_{t'}$ in the 550–600 GeV region is under less pressure than the one for $m_{b'}, m_{t'}$ in the 600–750 GeV region. Our analysis clearly shows that, even after imposing all collider and electroweak precision constraints, some regions of the parameter space survive, giving support to the hypothesis that the fourth generation is ruled out in models with a single doublet Higgs only, and survives when one adds a triplet Higgs field. We also find that a wide range of mass values for the sequential fourth generation quarks (∼550–750 GeV) is allowed by data provided that a very light (single- and double-) charged scalar spectrum with mass less than 210 GeV exists. These mass ranges are well within the reach of the current LHC experiment and would be a perfect testing ground to search for the complementary signals of the scalars and fermions of HTM4 in the current data and future runs.

ACKNOWLEDGEMENTS

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APPENDIX A: HIGGS DECAYS IN HTM4

In the HTM, in addition to the SM-like Higgs boson h, there are doubly charged scalar boson $H^{\pm\pm}$, singly charged scalar bosons $H^{\pm}$, one neutral CP-even scalar boson H

and one CP-odd scalar A. Their decays are affected by the presence of the fourth generation. Below, we present the decays for the neutral CP-even Higgs bosons h and H, as these are most likely to be candidates for the state at ~125 GeV observed at the LHC.

1. Decay rates of h

In the limit $v_\Delta \to 0$ the triplet Higgs decouples from the potential, and h is the SM-like Higgs boson.

(i) Tree-level decay rates for $h \to f\bar{f}, WW, WW^*, ZZ, ZZ^*$ The decay rates for h can be expressed as

$$\Gamma(h \to f\bar{f}) = G_F \frac{m_h m_f^2}{4\sqrt{2}\pi} N_c^f \beta\left(\frac{m_f^2}{m_h^2}\right)^3 \cos^2\alpha, \quad \text{(A1)}$$

$$\Gamma(h \to \nu\nu) = \Gamma(h \to \nu^c\bar{\nu}) + \Gamma(h \to \bar{\nu}^c\nu) = \sum_{i,j=1}^{4} S_{ij}|h_{ij}|^2 \frac{m_h}{4\pi}\sin^2\alpha, \quad \text{(A2)}$$

where $\beta(x) = \sqrt{1-4x}$. The decay rate of the Higgs boson h decaying into the gauge boson pair VV ($V = W$ or Z) is given by

$$\Gamma(h \to VV) = \frac{|\kappa_V(h)|^2 m_h^3}{128\pi m_V^4}\delta_V\left[1 - \frac{4m_V^2}{m_h^2} + \frac{12m_V^4}{m_h^4}\right] \times \beta\left(\frac{m_V^2}{m_h^2}\right), \quad \text{(A3)}$$

where $\delta_W = 2$ and $\delta_Z = 1$, and where $\kappa_V(h)$ are the couplings of the Higgs h with the vector bosons

$$\kappa_W(h) = \frac{ig^2}{2}(v_\phi \cos\alpha + 2v_\Delta \sin\alpha), \quad \text{(A4)}$$

$$\kappa_Z(h) = \frac{ig^2}{2\cos^2\theta_W}(v_\phi \cos\alpha + 4v_\Delta \sin\alpha). \quad \text{(A5)}$$

The decay rates for the three body decay modes are,

$$\Gamma(h \to VV^*) = \frac{3g_V^2|\kappa_V(h)|^2 m_h}{512\pi^3 m_V^2}\delta_{V'} F\left(\frac{m_V^2}{m_h^2}\right), \quad \text{(A6)}$$

where $\delta_{W'} = 1$ and $\delta_{Z'} = \frac{7}{12} - \frac{10}{9}\sin^2\theta_W + \frac{40}{27}\sin^4\theta_W$, and the function $F(x)$ is given as

$$F(x) = -|1-x|\left(\frac{47}{2}x - \frac{13}{2} + \frac{1}{x}\right) + 3(1 - 6x + 4x^2)|\log\sqrt{x}| + \frac{3(1-8x+20x^2)}{\sqrt{4x-1}}\arccos\left(\frac{3x-1}{2x^{3/2}}\right). \quad \text{(A7)}$$

Explicitly, inserting the couplings for the HTM4,

$$\Gamma(h \to W^+W^-) = \frac{g^4 m_h^3}{64\pi m_W^4}(v_\varphi \cos\alpha + 2v_\Delta \sin\alpha)^2 \times \left(\frac{1}{4} - \frac{m_W^2}{m_h^2} + \frac{3m_W^4}{m_h^4}\right)\beta\left(\frac{m_W^2}{m_h^2}\right), \quad \text{(A8)}$$

$$\Gamma(h \to ZZ) = \frac{g_Z^4 m_h^3}{128\pi m_Z^4}(v_\phi \cos\alpha + 4v_\Delta \sin\alpha)^2 \times \left(\frac{1}{4} - \frac{m_Z^2}{m_h^2} + \frac{3m_Z^4}{m_h^4}\right)\beta\left(\frac{m_Z^2}{m_h^2}\right), \quad \text{(A9)}$$

$$\Gamma(h \to WW^*) = \frac{3g^6 m_h}{2048\pi^3 m_W^2}(v_\phi \cos\alpha + 2v_\Delta \sin\alpha)^2 \times F\left(\frac{m_W^2}{m_h^2}\right), \quad \text{(A10)}$$

$$\Gamma(h \to ZZ^*) = \frac{g_Z^6 m_h}{8192\pi^3 m_Z^2}(v_\phi \cos\alpha + 4v_\Delta \sin\alpha)^2 \times \left(7 - \frac{40}{3}\sin^2\theta_W + \frac{160}{9}\sin^4\theta_W\right) \times F\left(\frac{m_Z^2}{m_h^2}\right). \quad \text{(A11)}$$

(ii) Loop induced decay rates for $h \to \gamma\gamma, gg$

$$\Gamma(h \to \gamma\gamma) = \frac{G_F \alpha^2 m_h^3}{128\sqrt{2}\pi^3}\left|\sum_f N_c^f Q_f^2 g_{hff} A_{1/2}^h(\tau_f) + g_{hWW}A_1^h(\tau_W) + g_{hH^\pm H^\mp}A_0^h(\tau_{H^\pm}) + 4g_{hH^{\pm\pm}H^{\mp\mp}}A_0^h(\tau_{H^{\pm\pm}})\right|^2. \quad \text{(A12)}$$

Here, α is the fine-structure constant, $N_c(=3)$ is the number of quark colors, Q_f is the electric charge of the fermion in the loop, and $\tau_i = m_h^2/4m_i^2$ ($i = f, W, H^\pm, H^{\pm\pm}$). The loop functions A_1 (for the W boson) and $A_{1/2}$ (for the fermions, f) are

$$A_{1/2}^h(\tau) = 2[\tau + (\tau-1)f(\tau)]\tau^{-2}, \quad \text{(A13)}$$

$$A_1^h(\tau) = -[2\tau^2 + 3\tau + 3(2\tau-1)f(\tau)]\tau^{-2}. \quad \text{(A14)}$$

For the contribution from the fermion loops, we will only keep the term with the t, b' and t' quarks, which are dominant. The loop function for $H^{\pm\pm}$ and $H^\pm$ is given by

$$A_0^h(\tau) = -[\tau - f(\tau)]\tau^{-2}, \quad \text{(A15)}$$

and the function $f(\tau)$ is given by

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & \tau \le 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & \tau > 1. \end{cases} \quad \text{(A16)}$$

Note that the contribution from the loop with $H^{\pm\pm}$ in Eq. (A12) is enhanced relative to the contribution from $H^{\pm}$ by a factor of four at the amplitude level. The couplings of h to the vector bosons and fermions relative to the values in the SM are as follows:

$$g_{hf\bar{f}} \equiv g_{ht\bar{t},hb'\bar{b}',ht'\bar{t}'} = \cos\alpha/\cos\beta_{\pm}, \quad \text{(A17)}$$

$$g_{hWW} = \cos\alpha + 2\sin\alpha v_{\Delta}/v_{\Phi}, \quad \text{(A18)}$$

$$g_{hZZ} = \cos\alpha + 4\sin\alpha v_{\Delta}/v_{\Phi}. \quad \text{(A19)}$$

The scalar trilinear couplings are parametrized as follows [36]:

$$g_{hH^{++}H^{--}} = \frac{m_W}{g m^2_{H^{\pm\pm}}}\{2\lambda_2 v_{\Delta}\sin\alpha + \lambda_1 v_{\Phi}\cos\alpha\} \quad \text{(A20)}$$

$$g_{hH^+H^-} = \frac{m_W}{2g m^2_{H^{\pm}}}\{[4v_{\Delta}(\lambda_2+\lambda_3)\cos^2\beta_{\pm} + 2v_{\Delta}\lambda_1\sin^2\beta_{\pm} - \sqrt{2}\lambda_4 v_{\Phi}\cos\beta_{\pm}\sin\beta_{\pm}]\sin\alpha + [\lambda v_{\Phi}\sin\beta_{\pm}{}^2 + (2\lambda_1+\lambda_4)v_{\Phi}\cos^2\beta_{\pm} + (4\mu - \sqrt{2}\lambda_4 v_{\Delta})\cos\beta_{\pm}\sin\beta_{\pm}]\cos\alpha\}. \quad \text{(A21)}$$

The loop induced decay process $h \to gg$ can be expressed by

$$\Gamma(h \to gg) = \frac{G_F\alpha_s^2 m_h^3}{64\sqrt{2}\pi^3}\cos^2\alpha\left|\sum_{q=1}^{4} A^h_{1/2}(\tau_q)\right|^2, \quad \text{(A22)}$$

where the loop functions are given by Eq. (A13), with $f(\tau)$ given by Eq. (A29).
Finally, the loop induced decay process $h \to Z\gamma$ is

$$\Gamma(h \to Z\gamma) = \frac{G_F\alpha_{\rm em}^2 m_h^3}{64\pi^4}\left(1 - \frac{m_Z^2}{m_h^2}\right)^3\left|\sum_f N_f^c Q_f g_{hf\bar{f}} A^h_{1/2}(\tau_f^{-1},\lambda_f^{-1}) + g_{hWW}A_1^h(\tau_W^{-1},\lambda_W^{-1}) + (g_{hH^+H^-})(g_{ZH^+H^-})A_0^h(\tau_{H^+}^{-1},\lambda_{H^+}^{-1}) + (g_{hH^{++}H^{--}})(g_{ZH^{++}H^{--}})A_0^h(\tau_{H^{++}}^{-1},\lambda_{H^{++}}^{-1})\right|^2, \quad \text{(A23)}$$

with $\tau_i = \frac{m_h^2}{4m_i^2}$, $\lambda_i = \frac{m_Z^2}{4m_i^2}$, and where the loop functions are given by

$$A^h_{1/2}(\tau^{-1},\lambda^{-1}) = \frac{2}{\cos\theta_W}(I_{3f} - 2Q_f\sin^2\theta_W) \times [I_1(\tau^{-1},\lambda^{-1}) - I_2(\tau^{-1},\lambda^{-1})], \quad \text{(A24)}$$

$$A_1^h(\tau^{-1},\lambda^{-1}) = \cos\theta_W\{4(3-\tan^2\theta_W)I_2(\tau^{-1},\lambda^{-1}) + [(1+2\tau)\tan^2\theta_W - (5+2\tau)]I_1(\tau^{-1},\lambda^{-1})\}, \quad \text{(A25)}$$

$$A_0^h(\tau_{H^c}^{-1},\lambda_{H^c}^{-1}) = I_1(\tau_{H^c}^{-1},\lambda_{H^c}^{-1}), \qquad c = +,++ \quad \text{(A26)}$$

$$I_1(\tau^{-1},\lambda^{-1}) = -\frac{1}{2(\tau-\lambda)} + \frac{1}{2(\tau-\lambda)^2}[f(\tau) - f(\lambda)] + \frac{\lambda}{(\tau-\lambda)^2}[g(\tau) - g(\lambda)] \quad \text{(A27)}$$

$$I_2(\tau^{-1},\lambda^{-1}) = \frac{1}{2(\tau-\lambda)}[f(\tau) - f(\lambda)]. \quad \text{(A28)}$$

In the equations above,

$$f(\tau) = \begin{cases} [\arcsin(\sqrt{\tau})]^2, & \text{if } \tau \le 1, \\ -\frac{1}{4}\left[\ln\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2, & \text{if } \tau > 1 \end{cases} \quad \text{(A29)}$$

and

$$g(\tau^{-1}) = \begin{cases} \sqrt{\tau-1}\left[\arcsin\left(\sqrt{\frac{1}{\tau}}\right)\right], & \text{if } \tau > 1, \\ \frac{\sqrt{1-\tau}}{2}\left[\ln\frac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}} - i\pi\right], & \text{if } \tau \le 1 \end{cases} \quad \text{(A30)}$$

and the couplings between the charged Higgs bosons and the Z are

$$g_{ZH^+H^-} = \frac{2}{\sin 2\theta_W}\frac{v_\Delta^2 \cos 2\theta_W - v_\phi^2 \sin\theta_W^2}{v_\phi^2 + 2v_\Delta^2} \quad \text{(A31)}$$

$$g_{ZH^{++}H^{--}} = \frac{2\cos 2\theta_W}{\sin 2\theta_W}. \quad \text{(A32)}$$

2. Decay rates of H

In the limit in which $\alpha \to 0$, the H Higgs boson is mostly triplet like; thus, this is the non-SM-like boson. The decays of H can be obtained from the corresponding formulas for h, with the substitution $\cos\alpha, \to -\sin\alpha,$, $\sin\alpha \to \cos\alpha$.

APPENDIX B: CONSTRAINTS ON THE HIGGS POTENTIAL

These have been thoroughly analyzed in [36], and we summarize their results briefly. The positivity requirement in the singly and doubly charged Higgs mass sectors, and choosing $v_\Delta > 0$, require

$$\mu > 0 \quad \text{(B1)}$$

$$\mu > \frac{\lambda_4 v_\Delta}{2\sqrt{2}} \quad \text{(B2)}$$

$$\mu > \frac{\lambda_4 v_\Delta}{\sqrt{2}} + \sqrt{2}\frac{\lambda_3 v_\Delta^3}{v_\Phi^2}, \quad \text{(B3)}$$

while for the requirement that the potential is bounded from below the complete set of conditions are

$$\lambda > 0 \;\&\; \lambda_2 + \lambda_3 > 0 \;\&\; \lambda_2 + \frac{\lambda_3}{2} > 0 \quad \text{(B4)}$$

$$\&\,\lambda_1 + \sqrt{\lambda(\lambda_2+\lambda_3)} > 0 \;\&\; \lambda_1 + \sqrt{\lambda\left(\lambda_2 + \frac{\lambda_3}{2}\right)} > 0 \quad \text{(B5)}$$

$$\&\,\lambda_1 + \lambda_4 + \sqrt{\lambda(\lambda_2+\lambda_3)} > 0$$
$$\&\,\lambda_1 + \lambda_4 + \sqrt{\lambda\left(\lambda_2 + \frac{\lambda_3}{2}\right)} > 0. \quad \text{(B6)}$$

Of these, the last expression in Eq. (B6) would restrict possible enhancements in the $h, H \to \gamma\gamma$ decay.

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