Double Higgs boson production as an exclusive probe for a sequential fourth generation with wrong-sign Yukawa couplings

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It has been shown that the data from the LHC do not rule out a chiral sequential fourth generation of fermions that obtain their masses through an identical mechanism as the other three generations do. However, this is possible only if the scalar sector of the Standard Model is suitably enhanced, like embedding it in a type-II two-Higgs doublet model. In this article, we try to show that double Higgs production (DHP) can unveil the existence of such a hidden fourth generation in a very efficient way. While the DHP cross section in the SM is quite small, it is significantly enhanced with a fourth generation. We perform a detailed analysis of the dependence of the DHP cross section on the model parameters and show that either a positive signal of DHP is seen in the early next run of the LHC or the model is ruled out.

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I. INTRODUCTION

More than seven years after the discovery of the Higgs boson at the LHC [1,2], the data have reached such a level of precision so as to vindicate the Standard Model (SM) [3] and to rule out a number of theories beyond the SM. The SM extended by a chiral fourth fermion generation (SM4)[4–10] constitutes such an example, as the quantum effects of SM4 on the production and decay of the Higgs boson do not decouple even in the heavy mass limit [11]. Especially, the production rate through gluon-gluon fusion, $gg \to h$, shoots up way beyond the experimental data.

However, it was recently shown in Ref. [12] that if the fourth generation is augmented by an extra Higgs doublet it is possible to hide such quantum effects completely, albeit in a certain case known as the wrong-sign limit [13–16]. The ambiguity in the determination of the sign of the down-type Yukawa couplings plays a crucial role in arranging cancellation among the amplitudes mediated by the fourth-generation fermions in the Higgs production and decay processes.[1] We achieve this conspiracy of Yukawa couplings in the framework of a type-II two-Higgs doublet model (2HDM). We would like to draw the attention of the readers to a couple of points here. First, there can be other extensions of the SM that can accommodate extra sequential generations of fermions in a similar way, and the type-II 2HDM is just an example. Second, such an arrangement saves the fourth generation only from the Higgs data. Other constraints and issues, like those coming from the oblique parameters, the stability of the potential, the scattering unitarity, and the inherent nonperturbative nature of the fourth-generation Yukawa couplings, must have to be taken care of separately, possibly by the introduction of other degrees of freedom and new dynamics associated with them.

When one talks about a second Higgs doublet, the data on the production and decay of the 125 GeV scalar resonance must be taken into account. In the framework of a 2HDM, which we will take to be type II for our discussion, this means that we must be close to the alignment limit [17–21]; i.e., the lighter CP-even neutral scalar h must be SM-like in its tree-level couplings to fermions and gauge bosons, modulo the phase of some of the couplings. The alignment limit and the wrong-sign limit are not the same, but they are close to each other when $\tan\beta(= v_2/v_1)$, the ratio of the vacuum expectation values (VEV) of the two Higgs doublets, is large and is related to the other parameters in a specific way.

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[1]The sign of the top Yukawa coupling with respect to the WWh coupling is fixed from the diphoton decay width of the Higgs boson.

Assuming that such a model—the SM extended by a heavy chiral fourth generation of leptons and quarks and another Higgs doublet of type-II variety, the complete package of which we will call xSM4—exists and somehow satisfies all the constraints mentioned in the previous paragraphs, one would like to ask the pertinent question of the possible direct signatures of the model. There can be indirect signals through flavor physics, coming from the mixing of the fourth generation with the other three, but that involves parameters that are *a priori* unknown. Similarly, direct production of heavy fermions also depends on unknown quantities like the masses of those fermions. The model can be tuned so as to make all such effects vanish.

There is, however, an interesting way to unveil the fourth generation, which we would like to focus on in this paper. This is the double Higgs production (DHP), $pp \to hh + X$. The DHP is the most important process to directly measure the Higgs self-coupling, and in the context of the LHC, this has been widely discussed in the literature [22–24]. DHP can proceed through various subprocesses: gluon-gluon fusion (ggF), vector boson fusion, Higgsstrahlung, or bremsstrahlung from the top; but ggF is by far the dominant process [25–27], contributing more than 90%. At $\sqrt{s} = 13$ TeV at the LHC, the total DHP cross section in the SM is 34.45 fb, of which 31.05 fb comes from ggF [25]. At $\sqrt{s} = 14$ TeV, the ggF share is 36.69 fb out of a total of 40.71 fb. In the ggF channel, there are two types of Feynman diagrams in the SM; one is represented by a top quark box, and the other is represented by a top quark triangle and a triple Higgs interaction. The DHP rate in the SM is suppressed because of a destructive interference between these two amplitudes. On the other hand, in xSM4, the new quarks play a pivotal role (as well as the new scalars), and the DHP cross section may receive an order-of-magnitude enhancement compared to the SM one, thanks to the large Yukawa couplings of the heavy fourth generation. We explain the mechanism in more detail in the next section.

We perform the analysis of DHP within the framework of xSM4 for $\sqrt{s} = 14$ TeV and show that (i) there should be a substantial enhancement of the DHP cross section, compared to that in the SM, which should either result in a positive signal very soon or rule the model out, and (ii) on top of the enhancement due to the fourth generation of fermions, there is a further contribution from the heavy neutral Higgs of 2HDM, the effect of which results in a severe constraint on the parameter space of 2HDM associated with xSM4.

At this point, one must point out the inherent limitations of such a study. Heavy chiral fermions that get their masses through Yukawa couplings naturally engender those couplings to be large and possibly nonperturbative. In this limit, the one-loop calculations have to be taken with a pinch of salt, as the higher-loop electroweak corrections may be even more significant [5,28]. However, the masses of the fourth-generation fermions can be tuned to arrange a cancellation [28], and consequently such effects can be diluted. In any case, such contributions should, in principle, raise the DHP cross section, if the higher-order terms add constructively to the leading-order ones. In other words, our bounds should turn out to be even stronger, or, in an extreme case, the model may already be ruled out. The ideal procedure would have been to use an effective theory, involving a $G^{a\mu\nu}G^{a}_{\mu\nu}\Phi^{\dagger}\Phi$ operator (i.e., involving two gluon and two Higgs fields) and integrating out the heavy fermion and scalar fields. However, this brings in an undetermined Wilson coefficient, the value of which must be ascertained by matching with the full theory. Thus, one has to evaluate the full ultraviolet-complete theory, as best as possible.

The paper is organized as follows. In Sec. II, we briefly recapitulate the DHP mechanism in the SM and also give a brief introduction to xSM4. The DHP in xSM4 is analyzed in Sec. III, with the results shown and discussed in Sec. IV. In the last section, we summarize our findings and conclude.

II. THEORY PRELUDE

A. Double Higgs production in the SM

The Feynman diagrams relevant for DHP in the SM ggF channel are shown in Fig. 1 (for all the diagrams, we refer the reader to the recent review, Ref. [25]). The subdominant contributions are neglected for our discussion.

The first diagram depicts an amplitude that proceeds through a quark box; this will be called a *box* (or $\square$) diagram. The second diagram similarly depicts an amplitude proceeding through a quark triangle and subsequent triple-Higgs interaction; this is called a *Delta* (or Δ) diagram. Note that there are many $\square$ and Δ diagrams

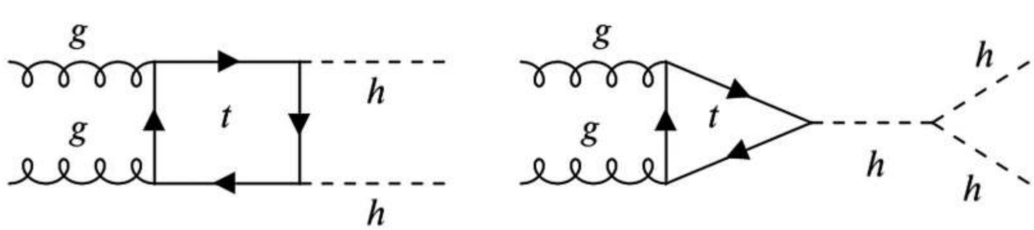

FIG. 1. Feynman diagrams for double Higgs production in the SM.

for different quarks; however, in the SM, one may neglect all other quarks except the top.

An interesting aspect of DHP in the SM is the destructive interference between the box and the Delta amplitudes. This results in an extremely small cross section, of about 36.69 fb at the 14 TeV LHC [25]. Searches for both resonant and nonresonant Higgs pair production have been performed in various channels by both the ATLAS and CMS experiments [29–39].

Let us display the relevant expressions for the DHP differential cross section via the ggF channel. The process under consideration is $g(p_1)g(p_2) \to h(p_3)h(p_4)$, where p_1 and p_2 are the incoming momenta, while p_3 and p_4 are the outgoing momenta, so that the Mandelstam variables are defined as

$$s=(p_1+p_2)^2, \quad t=(p_1-p_3)^2, \quad u=(p_1-p_4)^2. \tag{1}$$

To simplify the expressions, let us define some dimensionless quantities as

$$\mathcal{S}=\frac{s}{m_t^2}, \quad \mathcal{T}=\frac{t}{m_t^2}, \quad \mathcal{U}=\frac{u}{m_t^2}, \quad \rho=\frac{m_h^2}{m_t^2},$$
$$\mathcal{T}_1=\mathcal{T}-\rho, \quad \mathcal{U}_1=\mathcal{U}-\rho, \tag{2}$$

where m_t is the pole mass of the top quark and m_h is the Higgs boson mass. The spin and color averaged differential cross section for DHP is given by [40]

$$\frac{d\sigma_{gg\to hh}}{dt}=\frac{G_F^2\alpha_s^2}{256(2\pi)^3}[|(C_\Delta F_\Delta+C_\Box F_\Box)|^2+|C_\Box G_\Box|^2], \tag{3}$$

where G_F and α_s ($\equiv g_s^2/4\pi$) are the Fermi coupling and the strong coupling constants, respectively. The expressions for the F and G terms as well as the coefficients $C_\Box$ and C_Δ are given below:

$$C_{ij}=\int\frac{d^4q}{i\pi^2}\frac{1}{(q^2-m_t^2)[(q+p_i)^2-m_t^2][(q+p_i+p_j)^2-m_t^2]}, \tag{4}$$

$$D_{ijk}=\int\frac{d^4q}{i\pi^2}\frac{1}{(q^2-m_t^2)[(q+p_i)^2-m_t^2][(q+p_i+p_j)^2-m_t^2][(q+p_i+p_j+p_k)^2-m_t^2]}, \tag{5}$$

$$F_\Delta=\frac{2}{\mathcal{S}}\{2+(4-\mathcal{S})m_t^2C_{12}\}, \tag{6}$$

$$F_\Box=\frac{1}{\mathcal{S}^2}\{4\mathcal{S}+8m_t^2\mathcal{S}C_{12}-2m_t^4\mathcal{S}(\mathcal{S}+2\rho-8)(D_{123}+D_{213}+D_{132})$$
$$+2m_t^2(\rho-4)[\mathcal{T}_1(C_{13}+C_{24})+\mathcal{U}_1(C_{23}+C_{14})-m_t^2(\mathcal{T}\mathcal{U}-\rho^2)D_{132}]\}, \tag{7}$$

$$G_\Box=\frac{1}{\mathcal{S}(\mathcal{T}\mathcal{U}-\rho^2)}\{m_t^2(\mathcal{T}^2+\rho^2-8\mathcal{T})[\mathcal{S}C_{12}+\mathcal{T}_1(C_{13}+C_{24})-m_t^2\mathcal{S}\mathcal{T}D_{213}]$$
$$+m_t^2(\mathcal{U}^2+\rho^2-8\mathcal{U})[\mathcal{S}C_{12}+\mathcal{U}_1(C_{23}+C_{14})-m_t^2\mathcal{S}\mathcal{U}D_{123}]-m_t^2(\mathcal{T}^2+\mathcal{U}^2-2\rho^2)$$
$$\times(\mathcal{T}+\mathcal{U}-8)C_{34}-2m_t^4(\mathcal{T}+\mathcal{U}-8)(\mathcal{T}\mathcal{U}-\rho^2)(D_{123}+D_{213}+D_{132})\}. \tag{8}$$

The coefficients C_Δ and $C_\Box$ are given by

$$C_\Delta=\frac{3m_h^2}{s-m_h^2}g_t^h, \qquad C_\Box=(g_t^h)^2, \tag{9}$$

where $g_t^h\equiv\sqrt{2}m_t/v$ is the SM Yukawa coupling for the top quark. The cross section is calculated using FeynRules2.0 [41] and MadGraph5_aMC@NLO2.3.3 [42].

B. xSM4: A fourth chiral fermion generation in a type-II 2HDM

Here, we will summarize the essential points of Ref. [12]. One introduces a fourth chiral generation of fermions,

$$Q'=\begin{pmatrix}t'\\ b'\end{pmatrix}, \qquad L'=\begin{pmatrix}\nu'\\ \tau'\end{pmatrix}, \tag{10}$$

where all the components are massive enough to escape direct detection at the LHC and the masses are so arranged as to satisfy the bounds on both the oblique parameters S and T. Note that a full chiral generation, if completely degenerate, gives a constant contribution of $2/3\pi$ to the oblique S-parameter [4,43]. To cancel this large contribution, one has to introduce a mass splitting between the members of a doublet, and the splitting should go in opposite directions for the lepton and the quark doublets, to be in conformity with the T parameter. Keeping this in mind, we will display our results for two benchmark points that satisfy the bounds coming from the oblique parameters.

Within the ambit of the SM, there is no way to play with the sign of the Yukawa couplings; they must have the same sign as the masses. This is also necessary to maintain the tree-level unitarity [44]. However, if the scalar sector is extended to include another doublet, the wrong-sign limit, as shown below, may be achieved.

Let us focus on a type-II 2HDM with a CP-conserving scalar potential, where the two doublets are denoted by Φ_1 and Φ_2. We denote the VEVs of the two doublets by $v_1/\sqrt{2}$ and $v_2/\sqrt{2}$, respectively, with $v = \sqrt{v_1^2 + v_2^2}$ being the total electroweak VEV. The scalar spectrum has two CP-even neutral fields h and H, with h being the lighter one; one CP-odd neutral field A; and a pair of charged fields $H^\pm$. The physical fields h and H are obtained from the corresponding CP-even components of Φ_1 and Φ_2 by a rotation parametrized by the angle α. The fact that the lighter CP-even scalar, h, has couplings with the SM gauge bosons and fermions that are completely in conformity with the SM (within the experimental uncertainties) constrains the allowed parameter space to lie close to the alignment limit, defined by [17,18,21,45,46]

$$\cos(\beta - \alpha) \approx 0. \tag{11}$$

On the other hand, to cancel the fourth-generation contributions completely from the $gg \to h$, $h \to \gamma\gamma$, and $h \to Z\gamma$ processes, one needs the correlation among the scaling of the Higgs couplings relative to the SM,

$$\frac{g_{VVh}}{g_{VVh}^{\rm SM}} = \frac{g_{uuh}}{g_{uuh}^{\rm SM}} = -\frac{g_{ddh}}{g_{ddh}^{\rm SM}} = -\frac{g_{\ell\ell h}}{g_{\ell\ell h}^{\rm SM}} = 1, \tag{12}$$

where V, u, d, and ℓ denote, generically, the weak gauge bosons, the up-type quarks, the down-type quarks, and the charged leptons, respectively. Such a "wrong-sign limit" can be obtained, within the framework of a type-II 2HDM, as [13–16]

$$\cos(\beta - \alpha) = 2\cot\beta, \quad \text{with} \quad \tan\beta \gg 2. \tag{13}$$

This makes the down-type Yukawa couplings with the nonstandard neutral scalar large (they are enhanced by $\tan\beta$), and for b', the coupling can easily be nonperturbative. While this is an important issue that has recently been emphasized in Ref. [47], we will keep our analysis confined to not-too-large values of $\tan\beta$, with a tacit understanding that any other issues with the stability of xSM4 are somehow taken care of, most probably by some other dynamics.

III. DOUBLE HIGGS PRODUCTION IN xSM4

Now that we have outlined how, in the wrong-sign limit, the type-II 2HDM can hide a chiral fourth generation in processes like $gg \to h$, $h \to \gamma\gamma$, $h \to Z\gamma$, let us concentrate on DHP: $gg \to hh$. The relevant Feynman diagrams can be generalized from those shown in Fig. 1, by replacing t with $t/b'/t'$ and the scalar propagator with h/H. We follow the same path as outlined in Sec. II A, except that (i) all the three heavy quarks, viz., t, t', and b', will be taken into account[2] and, (ii) apart from the h-mediated Δ diagrams, the H-mediated diagrams have to be considered, too.

TABLE I. The benchmark points BP1 and BP2, with all masses in GeV.

Benchmark	$m_{t'}$	$m_{b'}$	$m_{\tau'}$	$m_{\nu'}$	m_H	$m_H^\pm$	m_A
BP1	1430	1380	1380	495	1010	1900	2800
BP2	1430	1380	1380	500	2160	2650	4050

It is quite obvious that the cancellation of the fourth-generation contributions due to the wrong-sign dynamics will no longer occur for the box diagrams, as the relevant Yukawa couplings appear twice in the box diagrams. Thus, we expect a large enhancement of the DHP rate over the SM prediction. Before we go into that, the Δ diagrams merit a comment. Note that there is a second Δ diagram, even in the absence of the fourth generation, that is mediated by the heavy scalar H. The Hhh coupling plays an important role in the DHP process, as we will see later. However, in the exact alignment limit $\cos(\beta - \alpha) = 0$, the Hhh coupling vanishes.

With more such $\Box$ and Δ diagrams in xSM4, expressions for whose amplitudes are analogous to what we have shown before, the differential cross section for DHP can be written as

$$\frac{d\sigma_{gg\to hh}}{dt} = \frac{G_F^2 \alpha_s^2}{256(2\pi)^3}\left[\left|\sum\nolimits_{t,t',b'}\{(C_\Delta^h + C_\Delta^H)F_\Delta + C_\Box F_\Box\}\right|^2 + \left|\sum\nolimits_{t,t',b'} C_\Box G_\Box\right|^2\right]. \tag{14}$$

While $F_\Box$, F_Δ, and $G_\Box$ have expressions analogous to those in Eq. (8), with the suitable replacement of the quark label, the expressions for C_Δ and $C_\Box$ are

$$C_\Delta^h = \lambda_{hhh}\frac{v}{s - m_h^2} g_q^h, \quad C_\Delta^H = \lambda_{hhH}\frac{v}{s - m_H^2} g_q^H, \quad C_\Box = (g_q^h)^2, \tag{15}$$

where $g_q^{h/H}$ $(q = t, t', b')$ are given by

[2]We refrain ourselves from going to such high values of $\tan\beta$ where the b-contributions become relevant and the Yukawa coupling for b' becomes badly nonperturbative.

$$g_q^h = \frac{\sqrt{2}m_q}{v}\frac{\cos\alpha}{\sin\beta} \quad \text{for } q = t, t', \qquad g_q^h = -\frac{\sqrt{2}m_q}{v}\frac{\sin\alpha}{\cos\beta} \quad \text{for } q = b',$$
$$g_q^H = \frac{\sqrt{2}m_q}{v}\frac{\sin\alpha}{\sin\beta} \quad \text{for } q = t, t', \qquad g_q^H = \frac{\sqrt{2}m_q}{v}\frac{\cos\alpha}{\cos\beta} \quad \text{for } q = b', \tag{16}$$

and

$$\lambda_{hhh} = -\frac{3}{v\sin 2\beta}[\lambda_S v^2\cos(\alpha+\beta)\cos^2(\beta-\alpha) - m_h^2\{2\cos(\alpha+\beta) + \sin 2\alpha\sin(\beta-\alpha)\}], \tag{17}$$

$$\lambda_{hhH} = \frac{\cos(\beta-\alpha)}{v\sin 2\beta}\left[\sin 2\alpha(2m_h^2 + m_H^2) - \frac{1}{2}\lambda_S v^2(3\sin 2\alpha - \sin 2\beta)\right], \tag{18}$$

where the dimensionless quantity [20,48],

$$\lambda_S = \frac{2}{v^2}\frac{m_{12}^2}{\sin\beta\cos\beta}, \tag{19}$$

is used as a convenient parametrization for the soft $\mathbb{Z}_2$ breaking effect in the scalar potential. For our analysis, we will use the dimensionless coupling $\tilde{g}$, defined as

$$\tilde{g}_{hhh(H)} \equiv \frac{v}{3m_h^2}\lambda_{hhh(H)}. \tag{20}$$

Note that, once the wrong-sign limit of Eq. (13) is imposed and some benchmark values for the nonstandard masses are chosen, the trilinear couplings are controlled by $\tan\beta$ and λ_S. Therefore, λ_S can be tuned properly to adjust the value of $\tilde{g}_{hhH}$, and consequently, the DHP cross section can be treated as a function of $\tilde{g}_{hhH}$ and $\tan\beta$.

We will start with the limit $\tilde{g}_{hhH} = 0$, i.e., when the H-mediated diagram is absent. In this case, the box amplitude is expected to pick up a factor of 3 compared to that in the SM due to the presence of three heavy quarks, t, t', and b'. Also notice that the t' and b' contributions should cancel each other in the h-mediated triangle amplitudes because of the wrong-sign limit. Since the t-mediated triangle amplitude is also subdominant, this will lead to an enhancement of the DHP cross section approximately by a factor of 9, which can be sensed in the near future.

Next, we slowly switch on $\tilde{g}_{hhH}$ till the experimental upper bound on the DHP cross section $\approx$ 330–340 fb [49] is reached. Thus, for a given $\tan\beta$, one obtains an upper bound on $|\tilde{g}_{hhH}|$, although the exact number depends on the sign of this coupling. We display our results for two different benchmark points as shown in Table I. The masses specify all the relevant parameters, like the scaled Yukawa couplings $g_q^{h(H)}$, as $\tan\beta$ is treated as a variable parameter and the mixing angle α is obtained from Eq. (13). The soft breaking term λ_S is adjusted so that that the DHP cross section reaches the experimental limit at $\tan\beta = 20$. The masses of the heavy scalars and the fourth-generation fermions are chosen in such a way as to evade the direct detection limits [50,51]. We have explicitly checked that the oblique parameters S and T for both the benchmark points as shown in Table I, coming from fermionic [4,43] and scalar [52,53] loops, are within the experimental limits given by [54]

$$\Delta S = 0.05 \pm 0.10, \qquad \Delta T = 0.08 \pm 0.12. \tag{21}$$

While computing the H-mediated Δ diagrams, we have taken into account the finite, possibly non-negligible widths of the nonstandard scalars in our numerical codes. However, the H-mediated diagrams are not the primary reasons behind the enhancement in the DHP cross section. Therefore, our main result will not crucially depend on the effects arising due to the finite width of the nonstandard scalars.

IV. OBSERVABILITY OF xSM4

As explained in the previous section, in the wrong-sign limit, $\tilde{g}_{hhH}$ and $\tan\beta$ are the only free parameters relevant for our study once we decide to stick to the chosen benchmark points. Keeping in mind that the up- and down-type Yukawa couplings in the fourth generation do not become badly nonperturbative and at the same time to be consistent with the wrong-sign limit, we confine ourselves to the range $3 < \tan\beta < 20$.

First, let us focus on the case in which the fourth-generation is present but the H-mediated diagrams are absent, possibly because of a vanishing $\tilde{g}_{hhH}$. The t'- and b'-mediated box amplitudes, on top of the t-mediated one, enhances the SM DHP cross section approximately by a factor of 9, as can be seen from Fig. 2, to about 300 fb. This is in contrast to the enhancement of the DHP cross section [55] due to resonant production of heavy scalars and their subsequent decays. The nature of the plot near $\tan\beta = 3$ can be explained by the imperfect cancellation of the t' and b' amplitudes in the h-mediated triangle diagrams. However, the cross section stabilizes for moderate or large

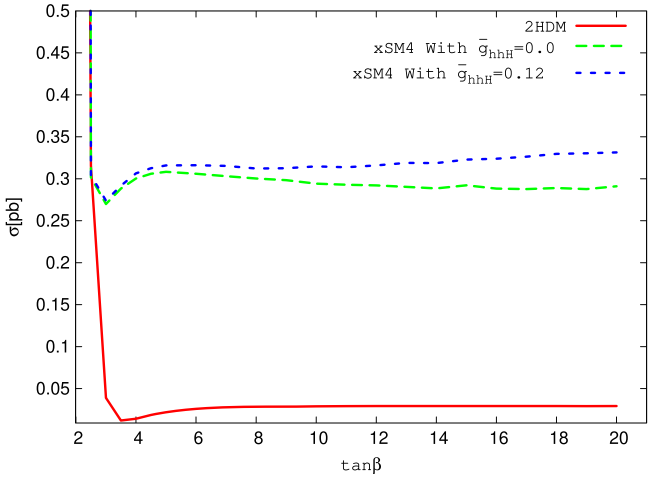

FIG. 2. Double Higgs production cross section as a function of $\tan\beta$ for benchmark point 1.

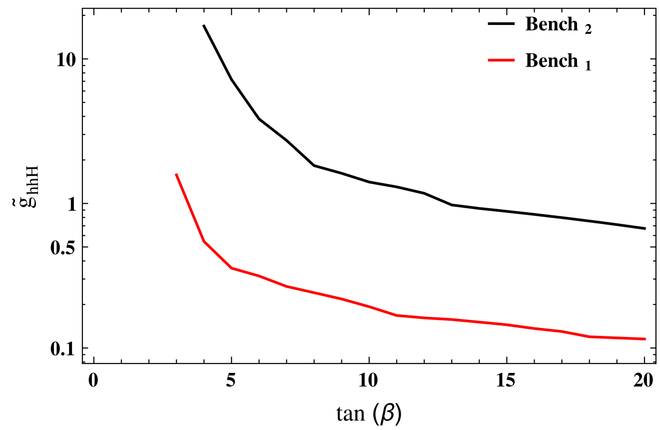

FIG. 3. The upper limit of $\tilde{g}_{hhH}$ as a function of $\tan\beta$ for the two benchmark points, with $M_H = 1.01$ and 2.16 TeV, respectively. Regions above the respective lines are ruled out from the experimental limit on DHP.

values of $\tan\beta$, i.e., when we are very close to the wrong-sign limit. We have checked that the nature of the plot remains the same for both the benchmark points. The region near $\tan\beta = 3$ is very close to, but not, the actual minimum of DHP cross section in xSM4. This will be discussed later in this section.

The Higgs pair going to the bb$\tau\tau$ channel is also studied in Ref. [56]. Among the all the probable final states into which h can decay, $hh \to b\bar{b}\tau^+\tau^-$ appears to be one of the most promising channels, because of a relatively small background compared to other final states [36]. The branching fraction of $hh \to b\bar{b}\tau^+\tau^-$ is about 7.3% [25] with the b-tagging efficiency being about 75% and the τ-tagging efficiency being about 50% [57–59]. So, the number of events observed, with an integrated luminosity of 300 fb^{-1}, is about 1200, which should be more than ample to verify the existence of the fourth generation in the wrong-sign limit.[3]

The upper limit of the DHP cross section, viz., 330–340 fb [49], is still somewhat above the enhancement caused by the fourth generation alone. This gap may be bridged with the contribution coming from the H-mediated Δ diagrams. This, in turn, sets an upper limit on $\tilde{g}_{hhH}$ as a function of $\tan\beta$, as displayed in Fig. 3. The limits are shown for the two benchmark points BP1 and BP2, with $m_H = 1.01$ and 2.16 TeV, respectively. The limits become stronger with increasing $\tan\beta$, as the Δ diagram with b' starts to become more significant. Obviously, the H-propagator suppression ensures that the limit on $\tilde{g}_{hhH}$ is higher for BP2 than for BP1; for $\tan\beta = 20$, the coupling can be as large as 0.67 for BP2 but only 0.12 for BP1.

The coupling $\tilde{g}_{hhH}$ can, in principle, be negative, too. Note that the sign is important only for the interference

[3]The present upper limit in this channel is 12.7 and 31.4 times that of the SM, by ATLAS and CMS collaborations, respectively.

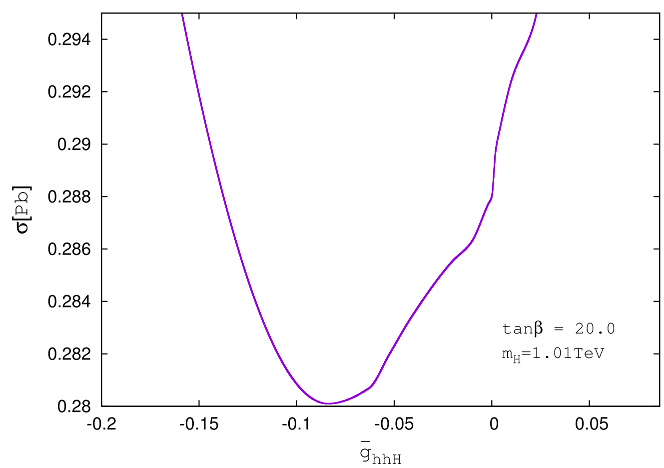

FIG. 4. Variation of the DHP cross section with $\tilde{g}_{Hhh}$, for $\tan\beta = 20$.

terms in the squared amplitude. It is easy to check that there is a destructive interference between the b'-mediated Δ diagram with the other $\square$ amplitudes, which slightly lowers the DHP cross section from its $\tilde{g}_{hhH} = 0$ value. This is shown in Fig. 4, in which we have taken $\tan\beta = 20$ to maximize the decrease. The lowest point of the curve is actually the minimum possible DHP cross section (approximately 280 fb) in xSM4, not the one that one gets from the absence of the Δ diagrams. However, the difference is negligible compared to the theoretical uncertainties in the estimation of the cross section as well as the experimental uncertainties, and even more so when the higher-order effects are neglected, as mentioned in the Introduction.

V. CONCLUSION

It is well known that DHP is the only process that can realistically probe the Higgs self-coupling λ. In a 2HDM that conserves CP, and that has all other scalars at 1 TeV or more except the 125 GeV Higgs, the DHP can probe the couplings λ_{hhh} and λ_{hhH}, the latter in the large $\tan\beta$ limit that enhances the bottom quark coupling to H.

In this paper, we show how DHP may be used to unveil a fourth chiral generation that can possibly hide in the LHC data. Such a model has been discussed in detail in Ref. [12], which is embedded in a type-II 2HDM. Even in the small parameter space where the fourth-generation effects on Higgs production and decay cancel out, the double Higgs production can successfully probe the model. This is because no such cancellation mechanism works for DHP, and one expects about an order-of-magnitude enhancement over the SM prediction, which is close to the experimental limit. Thus, the LHC at 14 TeV has an excellent chance to discover the signal, or rule the model out.

DHP limits can also rule out a significant portion of the parameter space of the scalar sector of xSM4, with the constraints on λ_{hhH} and $\tan\beta$. Admittedly, this also depends on other parameters of the potential, so one needs to look for other corroborative signals.

The analysis was performed at the leading order, and the higher-order effects may be extremely important, particularly when the Yukawa couplings are nonperturbative. Including such corrections should increase the DHP cross section and hence tighten our limits, or in an extreme case may rule out the model already. Even with this caveat, this remains an important and interesting channel to probe.

Before we conclude, it should be reemphasized that probing the trilinear self-coupling of the Higgs bosons has, so far, been the driving motivation behind the search for DHP at the LHC. Quite evidently, we can add to the stimulus for di-Higgs searches by making this channel a sensitive tool for unveiling certain new physics models. In this paper, we have done precisely that. We have shown that an extra sequential generation of fermions with wrong-sign Yukawa couplings, which can remain completely hidden in single Higgs production and decay [12], can potentially reveal themselves exclusively in the di-Higgs searches. To our knowledge, this is the first time such an exclusivity for the DHP in the context of new physics searches has been pointed out explicitly. Therefore, we hope that our current study will encourage our experimental colleagues to view the di-Higgs searches in a new light.

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