# Top quark flavor-changing neutral-current decay to a 125 GeV Higgs boson in the littlest Higgs model with $T$ parity

Bingfang Yang,[1,2,3] Ning Liu,[1,4,*] and Jinzhong Han[5]

[1]College of Physics & Electronic Engineering, Henan Normal University, Xinxiang 453007, China
[2]School of Materials Science and Engineering, Henan Polytechnic University, Jiaozuo 454000, China
[3]Basic Teaching Department, Jiaozuo University, Jiaozuo 454000, China
[4]ARC Centre of Excellence for Particle Physics at the Terascale, School of Physics, The University of Sydney, New South Wales 2006, Australia
[5]Department of Physics and Electronic Engineering, Zhoukou Normal University, Henan 466001, China

(Received 7 September 2013; revised manuscript received 20 November 2013; published 18 February 2014)

Motivated by the current observation of a 125 GeV Higgs boson, we calculate $t \to cH$ and $t \to cg(\gamma)H$ in the unitary gauge in the littlest Higgs model with $T$ parity (LHT). Because of the large contribution from the new mirror fermions, we find that the branching ratios of $t \to cH$ and $t \to cgH$ can be greatly enhanced in the LHT model and maximally reach $O(10^{-5})$ in the allowed parameter space. When the mirror fermion mass $M_3 > 2(1.5)$ TeV and the cutoff scale $f = 500$ GeV, the process of $pp \to t\bar{t} \to 3b + c + \ell + E_T^{\rm miss}$ can reach $3\sigma$ $(5\sigma)$ sensitivity at 8(14) TeV LHC with luminosity $\mathcal{L} = 20(300)$ fb$^{-1}$.

DOI: 10.1103/PhysRevD.89.034020 PACS numbers: 14.65.Ha, 12.15.Lk, 12.60.-i

## I. INTRODUCTION

In the light of the discovery of a Standard Model (SM)-like Higgs boson and the null results of new physics at the LHC, the electroweak hierarchy problem is highlighted much more than ever before. As the heaviest known elementary particle, the top quark has a strong correlation with the hierarchy problem and can be identified as a smoking gun of the TeV-scale physics.

In the SM, the top quark flavor-changing neutral-current (FCNC) processes are highly suppressed by the Glashow-Iliopoulos-Maiani (GIM) mechanism [1]. It indicates that any observation of these processes will be a signal beyond the SM [2–6]. Since weakly constrained FCNC couplings between the second- and the third-generation up-type quarks are usually predicted in some new physics models, the two-body FCNC decays $t \to cX (X = g, \gamma, Z, H)$ can be greatly enhanced, such as in the minimal supersymmetric Standard Model (MSSM) with branching ratio ${\rm Br}(t \to cH) \sim 10^{-5}$ [7], $R$-parity violating supersymmetry (SUSY) with branching fraction ${\rm Br}(t \to cH) \sim 10^{-6}$ [8], the two-Higgs doublet model (2HDM) with branching ratio ${\rm Br}(t \to cH) \sim 10^{-3}$ [9], and so on. The NLO QCD corrections to $t \to qH (q = u, c)$ in a model-independent method have been studied in Ref. [10]. In addition, three-body FCNC decays of the top quark were also found to be a sensitive probe of new physics, such as $t \to cX_1X_2 (X = g, \gamma, Z, H)$ [11–13]. Very recently, ATLAS Collaboration has measured the top quark decays $t \to cH$ with $H \to \gamma\gamma$ and set the upper limit on the $tcH$ coupling as 0.17 [14].

As an extension of the SM, the littlest Higgs model with $T$ parity (LHT) [15] can successfully solve the electroweak hierarchy problem by constructing the Higgs as a pseudo-Goldstone boson. Meanwhile, the discrete symmetry $T$ parity in this model also forbids the tree-level contributions from the heavy gauge bosons; thus, the LHT can safely avoid the constraint from the electroweak precision observables (EPO) that occurs in the littlest Higgs (LH) model [16]. For the top quark sector in the LHT, the top quark can interact with new $T$-odd gauge bosons and $T$-odd fermions, which may produce large contributions to the top quark FCNC processes [17]. Similar effects have been studied in the rare decays of the $K/B$ meson [18], the Higgs boson [19], and the $Z$ boson [19].

It should be mentioned that the searches for the LHT particles at the LHC can provide direct evidence of the LHT model or give strong constraints on the LHT parameter space. However, the results usually depend on the assumption of the specific mass spectrum and the branching ratios. For example, the $T$-odd top partner $(T^-)$ pair production has been explored through $pp \to T^-T^- \to A_H t A_H \bar{t}$ at 7 TeV LHC [20]. In the analysis, a large mass splitting between the $A_H$ and $T^-$ is required to produce the hard missing energy to suppress the top pair background. But in a general LHT model, the mass of $T^-$ can be close to the mass of $(A_H + t)$ or $A_H$ so that the adopted strategy is not applicable. Similar things can also happen in the searches for other LHT new particles. So in these cases, the searches for the indirect LHT effects via loop corrections will be of great importance because of the model's

*wlln@mail.ustc.edu.cn

 

weak dependence on the kinematics information. In particular, the processes with low SM backgrounds, such as top quark FCNC decays, will be helpful for testing the LHT model.

In this work, we calculate the top quark FCNC decays with Higgs interactions in unitary gauge in the LHT, that is, $t \to cH$ and $t \to cg(\gamma)H$. As a top quark factory, 14 TeV LHC has the power to detect the branching ratios of $t \to cH$ up to $\mathrm{Br} \sim O(10^{-6})$ for $\mathcal{L} = 30\ \mathrm{fb}^{-1}$ and $\mathrm{Br} \sim O(10^{-7})$ for $\mathcal{L} = 300\ \mathrm{fb}^{-1}$ [21]. So the study of these top FCNC processes can be used to test the LHT at the LHC. The paper is organized as follows. In Sec. II we recapitulate the LHT model related to our work. In Secs. III and IV we calculate the one-loop contributions of the LHT model to the $t \to cH$ and $t \to cg(\gamma)H$ in unitary gauge and present the numerical results. Finally, we give a short summary in Sec. V.

## II. A BRIEF REVIEW OF THE LHT MODEL

The LHT model is a nonlinear $\sigma$ model based on the coset space $SU(5)/SO(5)$, with the $SU(5)$ global symmetry broken by the vacuum expectation value (VEV) of a $5 \times 5$ symmetric tensor,

$$\Sigma_0 = \begin{pmatrix} \mathbf{0}_{2\times2} & 0 & \mathbf{1}_{2\times2} \\ 0 & 1 & 0 \\ \mathbf{1}_{2\times2} & 0 & \mathbf{0}_{2\times2} \end{pmatrix}. \tag{1}$$

The VEV of $\Sigma_0$ breaks the extended gauge group $[SU(2) \times U(1)]^2$ down to the SM electroweak $SU(2)_L \times U(1)_Y$, which leads to new heavy gauge bosons $W_H^{\pm}, Z_H, A_H$ with the masses given to lowest order in $v/f$ by

$$M_{W_H} = M_{Z_H} = gf\left(1 - \frac{v^2}{8f^2}\right), \qquad M_{A_H} = \frac{g'f}{\sqrt{5}}\left(1 - \frac{5v^2}{8f^2}\right). \tag{2}$$

Here $g$ and $g'$ are the SM $SU(2)$ and $U(1)$ gauge couplings, respectively.

When $T$ parity is implemented in the quark sector of the model, we require the existence of mirror partners with $T$-odd quantum number for each SM quark. We denote them by $u_H^i, d_H^i$, where $i$ $(i = 1, 2, 3)$ is the generation index. After electroweak symmetry breaking, a small mass splitting between $u_H^i$ and $d_H^i$ is induced, and the masses are given by

$$m_{d_H^i} = \sqrt{2}\kappa_i f, \qquad m_{u_H^i} = m_{d_H^i}\left(1 - \frac{v^2}{8f^2}\right), \tag{3}$$

where $\kappa_i$ are the diagonalized Yukawa couplings of the mirror quarks.

In order to stabilize the Higgs mass, an additional $T$-even heavy quark $T^+$ is introduced to cancel the large one-loop quadratic divergences caused by the top quark. But the implementation of $T$ parity requires a $T$-odd mirror partner $T^-$ with $T^+$. Their masses are given by

$$m_{T^+} = \frac{f}{v}\frac{m_t}{\sqrt{x_L(1-x_L)}}\left[1 + \frac{v^2}{f^2}\left(\frac{1}{3} - x_L(1-x_L)\right)\right], \tag{4}$$

$$m_{T^-} = \frac{f}{v}\frac{m_t}{\sqrt{x_L}}\left[1 + \frac{v^2}{f^2}\left(\frac{1}{3} - \frac{1}{2}x_L(1-x_L)\right)\right], \tag{5}$$

where $x_L$ is the mixing parameter between the SM top quark and its heavy partner $T^+$.

In the LHT model, the mirror quark Yukawa interaction is given by

$$\mathcal{L}_{\mathrm{mirror}} = -\kappa_{ij} f(\bar{\Psi}_2^i \xi + \bar{\Psi}_1^i \Sigma_0 \Omega \xi^\dagger \Omega)\Psi_R^j + \mathrm{H.c.} \tag{6}$$

A new flavor structure can come from the mirror fermions when the mass matrix $\sqrt{2}\kappa_{ij} f$ is diagonalized by two $U(3)$ matrices. One of the important ingredients of the mirror quark sector is the existence of two CKM-like unitary mixing matrices: $V_{Hu}, V_{Hd}$. These mirror mixing matrices parametrize flavor-changing interactions between SM quarks and mirror quarks that are mediated by the heavy gauge bosons $W_H^{\pm}, Z_H, A_H$.

Note that $V_{Hu}$ and $V_{Hd}$ are related through the SM CKM matrix:

$$V_{Hu}^\dagger V_{Hd} = V_{\mathrm{CKM}}. \tag{7}$$

We follow Ref. [22] to parametrize $V_{Hd}$ with three angles $\theta_{12}^d, \theta_{23}^d, \theta_{13}^d$ and three phases $\delta_{12}^d, \delta_{23}^d, \delta_{13}^d$,

$$V_{Hd} = \begin{pmatrix} c_{12}^d c_{13}^d & s_{12}^d c_{13}^d e^{-i\delta_{12}^d} & s_{13}^d e^{-i\delta_{13}^d} \\ -s_{12}^d c_{23}^d e^{i\delta_{12}^d} - c_{12}^d s_{23}^d e^{i(\delta_{13}^d - \delta_{23}^d)} & c_{12}^d c_{23}^d - s_{12}^d s_{23}^d s_{13}^d e^{i(\delta_{13}^d - \delta_{12}^d - \delta_{23}^d)} & s_{23}^d c_{13}^d e^{-i\delta_{23}^d} \\ s_{12}^d s_{23}^d e^{i(\delta_{12}^d + \delta_{23}^d)} - c_{12}^d c_{23}^d s_{13}^d e^{i\delta_{13}^d} & c_{12}^d s_{23}^d e^{i\delta_{23}^d} - s_{12}^d c_{23}^d s_{13}^d e^{i(\delta_{13}^d - \delta_{12}^d)} & c_{23}^d c_{13}^d \end{pmatrix}. \tag{8}$$

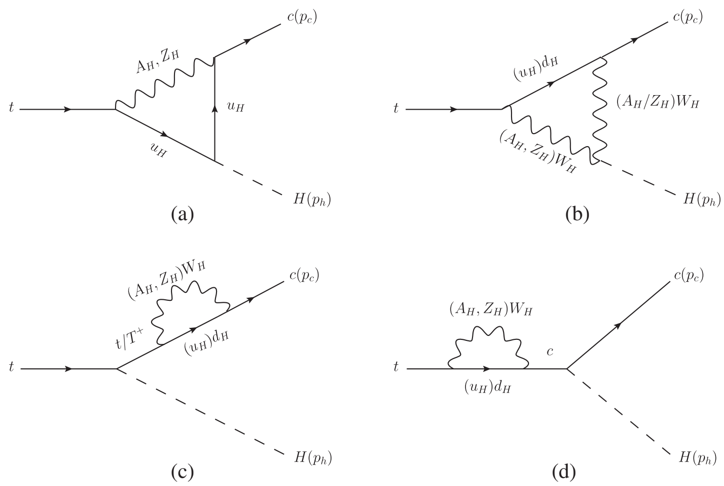

FIG. 1. Feynman diagrams of the LHT one-loop correction to $t \to cH$ in the unitary gauge.

In our calculation, for the matrices $V_{Hu}, V_{Hd}$, to aid comparisons with Ref. [23], we also follow Ref. [24] to consider the same scenarios as follows:

(i) Case I: $V_{Hd} = \mathbf{1}$,

(ii) Case II: $s_{12}^d = \frac{1}{\sqrt{2}}, s_{23}^d = 5 \times 10^{-5}, s_{13}^d = 4 \times 10^{-2}$, $\delta_{12}^d = \delta_{23}^d = \delta_{13}^d = 0$,

(iii) Case III: $s_{12}^d = 0.99, s_{23}^d = 2 \times 10^{-4}, s_{13}^d = 0.6$, $\delta_{12}^d = \delta_{23}^d = \delta_{13}^d = 0$.

## III. BRANCHING RATIO FOR $T \to CH$ IN THE LHT MODEL

In the LHT model, the relevant Feynman diagrams of the process $t \to cH$ in unitary gauge are shown in Fig. 1. We can see that there is no additional mixing between $T^+$ and the charm or up quark. This is different from the case in Ref. [25], where a small loop-induced coupling between a new vectorlike quark and a charm quark can occur and will be constrained by the low energy physics. We will not consider the higher order couplings between the scalar triplet $\Phi$ and top quark, and we neglect the high order $O(v^2/f^2)$ terms in the masses of new particles. The calculations of the loop diagrams are straightforward. We adopt the definitions of scalar one-loop integral functions in Ref. [26], and compose each loop diagram into some scalar loop functions [27]. We use the package LOOPTOOLS [28] to perform the numerical loop calculations. In the analytic calculations, we cancel the divergence that is independent on the mirror quark mass by the unitarity of the matrix $V_{Hu}$. We also note that the modified interactions of the up-type mirror fermions with the $Z$ boson in Ref. [29] can cancel similar divergences in the processes with down-type quarks or leptons as the external particles. However, we checked that such a modification cannot cancel the divergence in $t \to cH$, and there is still a so-called leftover divergence [23,30] as follows:

$$D = \frac{m_{u_H^i}^2}{f^2} (V_{Hu})_{i2}^* (V_{Hu})_{i3} \frac{i}{16\pi^2} \times \left[ -\frac{1}{80} + \left( \frac{x_L^2}{160} + \frac{3}{64} \right) \frac{v^2}{f^2} \right] \Delta, \tag{9}$$

where $\Delta = \frac{1}{\varepsilon} - \gamma_E + \ln 4\pi$. In Ref. [30], this so-called leftover divergence was understood as the sensitivity of the decay amplitudes to the ultraviolet completion of the LHT model. Following Ref. [30], we remove the divergent term $1/\varepsilon$ in the invariant amplitudes and take the renormalization scale $\mu = \Lambda$, with $\Lambda = 4\pi f$ being the cutoff scale of the LHT model.

In our numerical calculations, the SM parameters are taken as follows [31]:

$$G_F = 1.16637 \times 10^{-5}\ \mathrm{GeV}^{-2}, \qquad \sin^2\theta_W = 0.231, \qquad \alpha_e = 1/128, \qquad \alpha_s = 0.1076,$$
$$m_c = 1.27\ \mathrm{GeV}, \qquad M_Z = 91.1876\ \mathrm{GeV}, \qquad m_t = 173.5\ \mathrm{GeV}, \qquad m_h = 125\ \mathrm{GeV}. \tag{10}$$

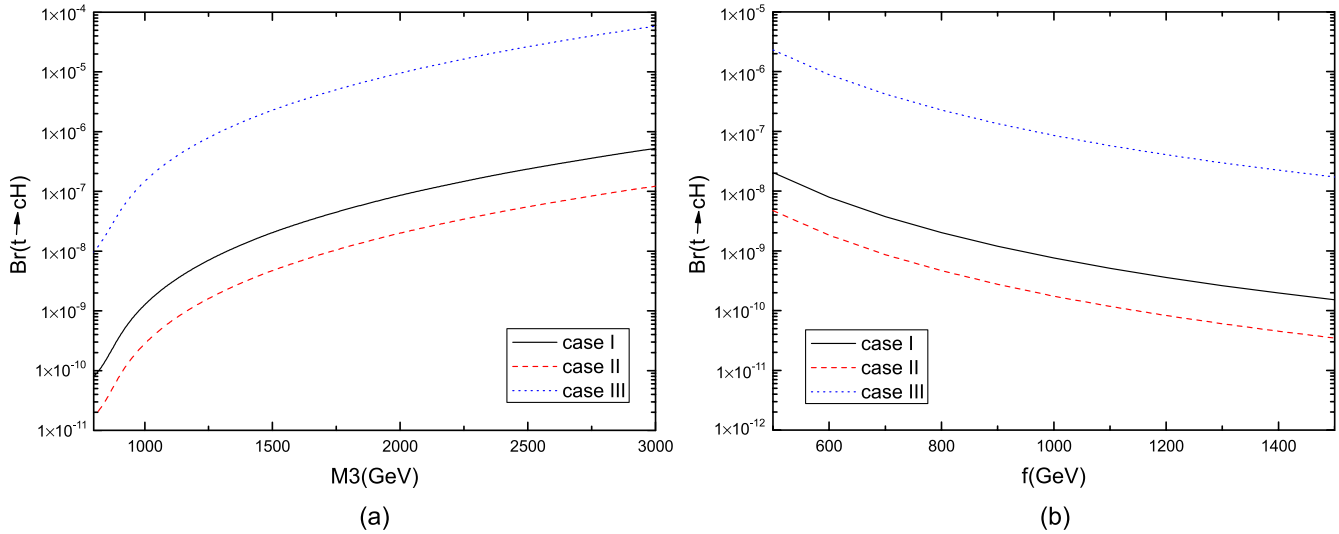


FIG. 2 (color online). Branching ratios of $t \to cH$ as a function of $M_3$ (a) and $f$ (b) in three cases, respectively.

The relevant LHT parameters in our calculation are the scale $f$, the mixing parameter $x_L$, the mirror quark masses, and the parameters in the matrices $V_{Hu}, V_{Hd}$. Considering the constraints in Refs. [32], the scale $f$ can be allowed to be as low as 500 GeV. For the mirror quark masses, it has been shown that the experimental bounds on four-fermi interactions require $m_{Hi} \leq 4.8 f^2/\mathrm{TeV}$ [33]; we get $m_{u_H^i} = m_{d_H^i}$ at $O(v/f)$ and further assume

$$m_{u_H^1} = m_{u_H^2} = m_{d_H^1} = m_{d_H^2} = M_{12}, \qquad m_{u_H^3} = m_{d_H^3} = M_3. \tag{11}$$

In Fig. 2(a), we show the dependence of the branching ratio of $t \to cH$ on the third-generation mirror quark mass $M_3$. We set the scale $f = 500$ GeV, the mixing parameter $x_L = 0.1$, and the first two-generation mirror quark masses $M_{12} = 750$ GeV. Because of the large departures from the SM caused by the mixing matrices in case III, we can also see that the branching ratio of $t \to cH$ in case III is much larger than in cases I and II, which can maximally reach $5.8 \times 10^{-5}$ in case III.

From Fig. 2(a), we can see that the branching ratio of $t \to cH$ increases as $M_3$ increases, which means that the decay rate is enhanced by the mass splitting between the three-generation mirror quarks. According to the analytic expression, we know the form factors of the loop-induced $tcH$ interaction, $F$, should take the following form:

$$F \propto \sum_{i=1}^{3} (V_{Hu}^{\dagger})_{ci} f(m_{Hi}) (V_{Hu})_{it}. \tag{12}$$

where $f(m_{Hi})$ is a universal function for three-generation mirror quarks, but its value depends on the mass of $i$th-generation mirror quark, $m_{Hi}$. Obviously, for the degeneracy of the three-generation mirror quarks, $F$ vanishes exactly due to the unitary of $V_{Hu}$, while for the degeneracy of the first two generations as discussed below, the factor behaves like $(V_{Hu}^{\dagger})_{c3}(f(m_{H3}) - f(m_H))(V_{Hu})_{3t}$, with $m_H$ being the common mass of the first two generations. The decay rate is enhanced by the mass splitting between the three-generation mirror quarks; since we set $M_1 = M_2 = M_{12}$, there is only one mass splitting $M_3 - M_{12}$, which increases with $M_3$ while keeping $M_{12}$ fixed. This agrees with the explanation in Ref. [17].

In Fig. 2(b), we show the dependence of the branching ratio of $t \to cH$ on the scale $f$. We set the mixing parameter $x_L = 0.1$, the first two-generation mirror quark masses $M_{12} = 1.5f$, and the third-generation mirror quark mass $M_3 = 3f$. We can see that the branching ratio decreases as the scale $f$ increases, which means that the correction of the LHT model decouples as the scale $f$ increases. Since the enhancement from the mass splitting of mirror fermions can balance the suppression of large scale $f$, we find that the branching ratio of $t \to cH$ decreases slowly when the scale $f$ becomes higher. From Fig. 2, we can see that the LHT model can enhance the branching ratios of $t \to cH$ as much as 9–10 orders of the one in the SM [34]. Similarly, in

TABLE I. Branching ratio for top quark decay $t \to cH$ in different models.

| | SM | QS | 2HDM | FC 2HDM | MSSM | $R$ SUSY | SUSY-QCD | LHT |
|---|---|---|---|---|---|---|---|---|
| $t \to cH$ | $3 \times 10^{-15}$ | $4.1 \times 10^{-5}$ | $1.5 \times 10^{-3}$ | $\sim 10^{-5}$ | $10^{-5}$ | $\sim 10^{-6}$ | $\sim 10^{-5}$ | $\sim 10^{-5}$ |

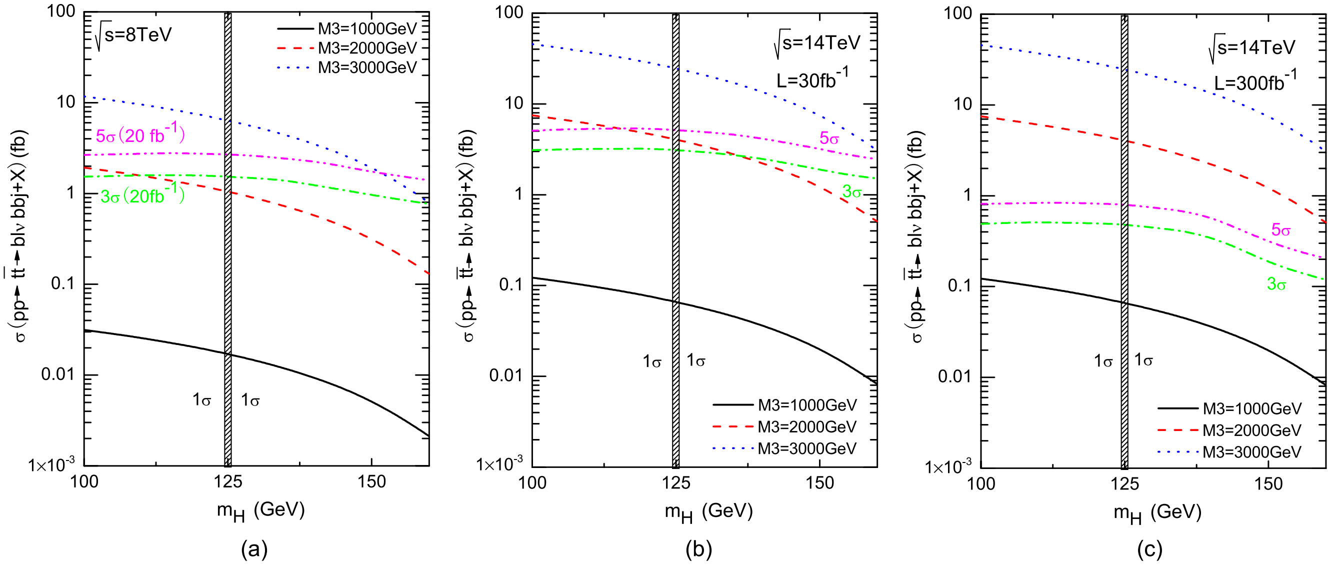


FIG. 3 (color online). The observability of $t \to cH$ for case III in the LHT model through the production of $pp \to t\bar{t} \to tcH \to bl\nu cb\bar{b} + X$ at the LHC with $\sqrt{s} = 8$ TeV and $\sqrt{s} = 14$ TeV. The shadow region is the $1\sigma$ combined range of the Higgs boson mass from Ref. [40].

some other new physics beyond the SM this branching ratio can also be enhanced by several orders of magnitude. For comparison, we summarize the FCNC decays $t \to cH$ in the LHT model and in other new physics models [8,35–38] in Table I.

At the LHC, the dominant background of the search for $t \to cH$ is the final state of $4j/3b\ell\nu$ coming from top quark pair production: $pp \to t\bar{t} \to b\ell^+\nu\bar{b}\,\bar{c}\,s + X$ or $pp \to t\bar{t} \to b\ell^-\bar{\nu}bc\bar{s} + X$, where a $c$ jet is misidentified as a $b$ jet. The mistagged probability of a $c$ jet as a $b$ jet is approximately 10% as reported by the ATLAS and CMS collaborations. In order to investigate the observability of $t \to cH$ for case III in the LHT model, we use the Monte Carlo simulation results in Ref. [39] and plot $3\sigma$ and $5\sigma$ contours of the hadronic cross sections $pp \to t\bar{t} \to b\ell\nu b\bar{b}j$ in Fig. 3 for $\sqrt{s} = 8, 14$ TeV. We use the next-to-leading order value of the $t\bar{t}$ production rate in the calculation. Since the branching ratio of $t \to cH$ is sensitive to the third-generation mirror quark mass, we take $M_3 =$ 1000 GeV, 2000 GeV, 3000 GeV for example, where we set the scale $f = 500$ GeV, the mixing parameter $x_L = 0.1$, and the first two-generation mirror quark masses $M_{12} = 750$ GeV.

On the left panel of Fig. 3, we can see that when $M_3 > 2.2$ TeV, $t \to cH$ can reach $3\sigma$ sensitivity at 8 TeV LHC with luminosity $\mathcal{L} = 20\ \text{fb}^{-1}$. But on the middle and

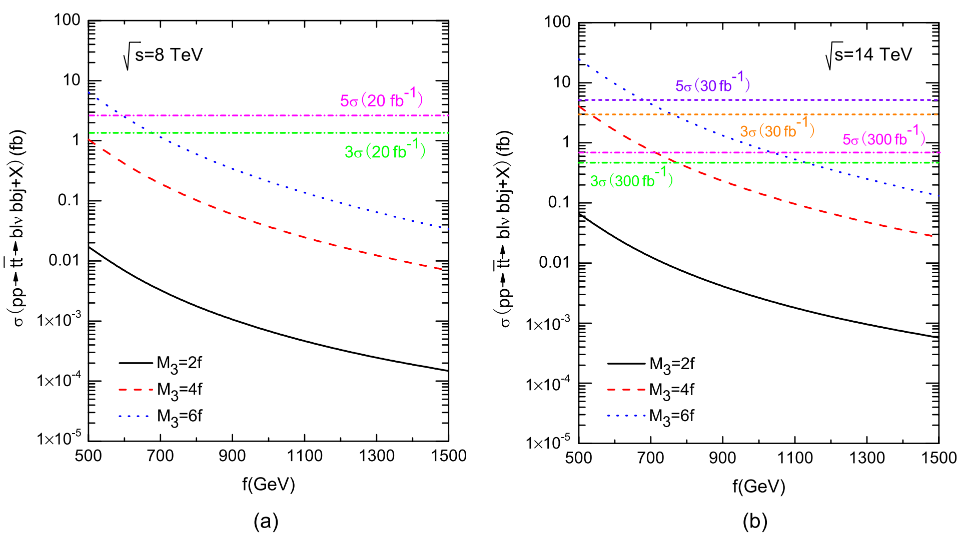


FIG. 4 (color online). The observability of $t \to cH$ as a function of the scale $f$ for three cases in the LHT model through the production of $pp \to t\bar{t} \to tcH \to bl\nu cb\bar{b} + X$ at the LHC with $\sqrt{s} = 8$ TeV and $\sqrt{s} = 14$ TeV.

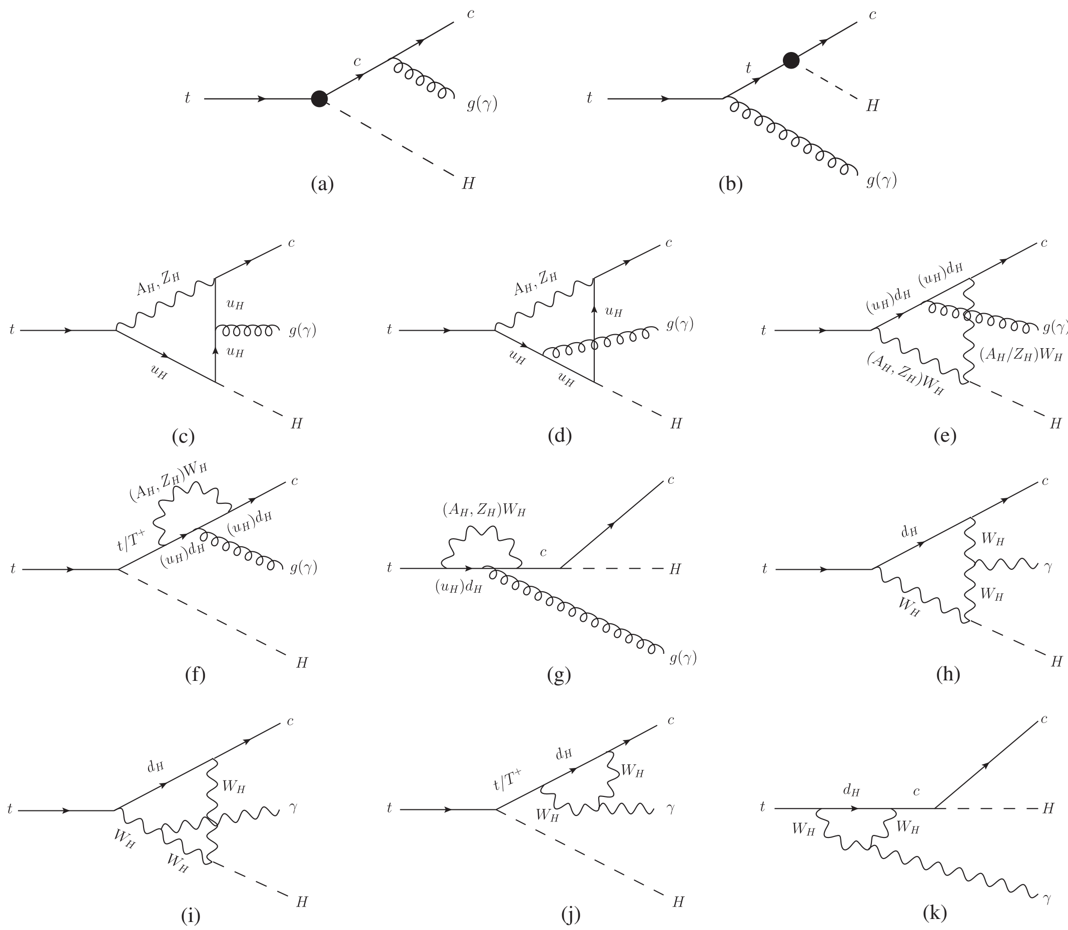


FIG. 5. Feynman diagrams of the LHT one-loop correction to $t \to cg(\gamma)H$ in the unitary gauge.

right panels of Fig. 3, we find that the 14 TeV LHC has the ability to probe a value of $M_3$ as low as 2.1(1.5) TeV at the $5\sigma$ level when $\mathcal{L} = 30(300)\ \mathrm{fb}^{-1}$. Therefore, we can infer that the precise measurement of $t\bar{t}$ production can give a strong constraint on the parameter space of the LHT model.

In Fig. 4, we show the observability of $t \to cH$ as a function of the scale $f$ for case III in the LHT model through the production of $pp \to t\bar{t} \to tcH \to bl\nu cb\bar{b} + X$ at the LHC with $\sqrt{s} = 8$ TeV and $\sqrt{s} = 14$ TeV. We use the Monte Carlo simulation results and the next-to-leading

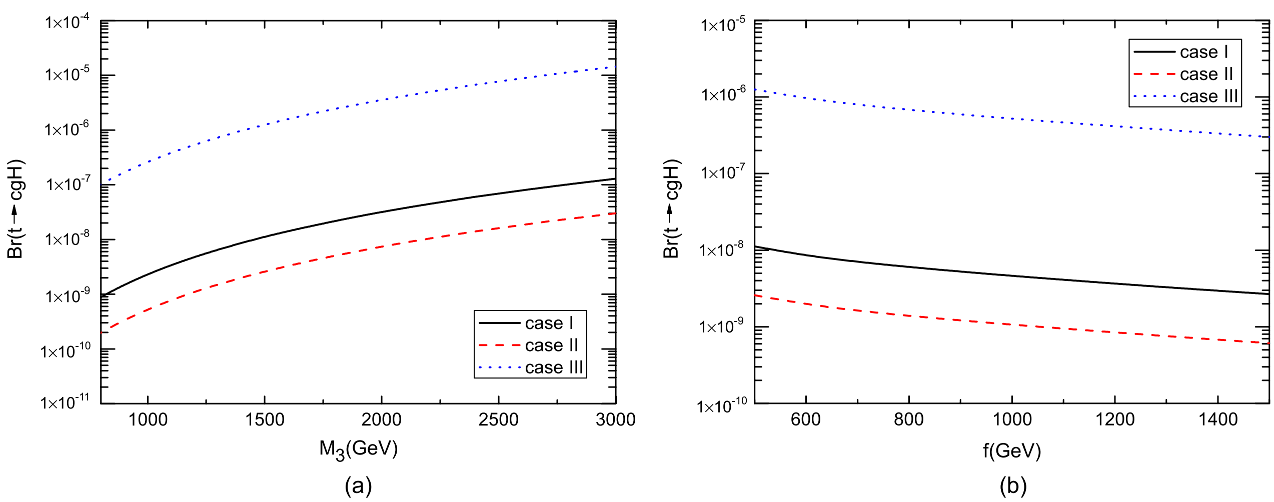


FIG. 6 (color online). Branching ratios for $t \to cgH$ as a function of $M_3$ (a) and $f$ (b) in three cases, respectively.

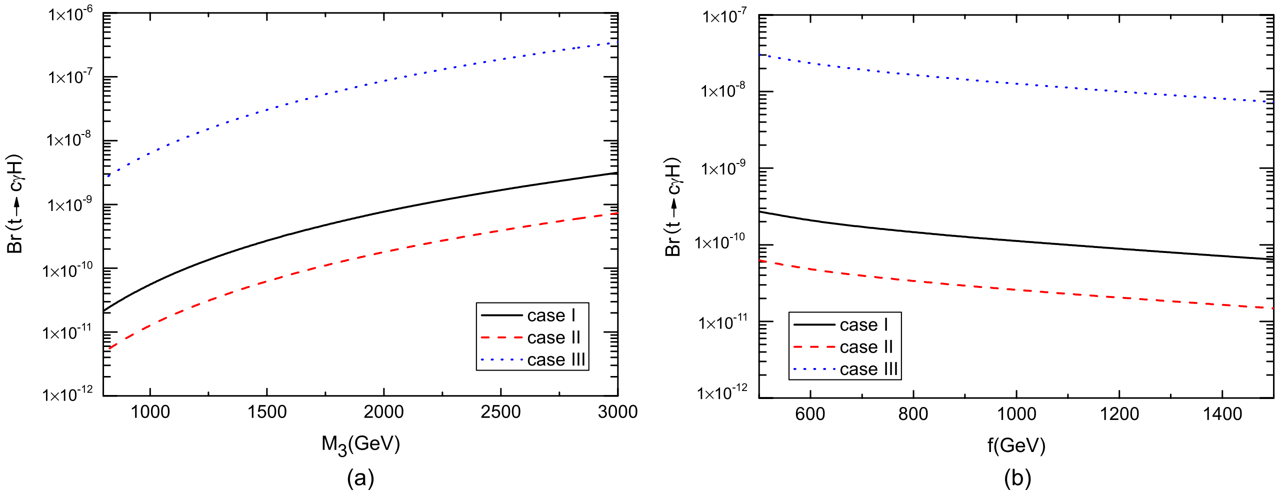


FIG. 7 (color online). Branching ratios for $t \to c\gamma H$ as a function of $M_3$ (a) and $f$ (b) in three cases, respectively.

order value of the $t\bar{t}$ production rate as above. The relevant parameters are taken as follows: $x_L = 0.1$, $M_{12} = 1.5f$, $m_h = 125$ GeV. Based on the same consideration, we take the third-generation mirror quark mass $M_3 = 2f$, $4f, 6f$, for example. We can see that the favorable observability comes from the region with the low $f$ and the large mass splitting $(M_3 - M_{12})$, which is consistent with the preceding analysis.

## IV. BRANCHING RATIO FOR $T \to CG(\gamma)H$ IN THE LHT MODEL

In this section, we calculate the branching ratio of $t \to cg(\gamma)H$ in the LHT model. These processes can also be considered as part of the next-to-leading order QCD (QED) corrections to $t \to cH$. The relevant Feynman diagrams of the process $t \to cg(\gamma)H$ in unitary gauge are shown in Fig. 5, where the black dots represent the loop-induced $tcH$ vertex, as shown in Fig. 1. In the numerical calculations, we take the same parameters and cases as the decay process $t \to cH$ and impose the kinematical cuts on the final massless states to avoid the singularity.

In Fig. 6(a), we show the dependence of the branching ratio of the $t \to cgH$ decay process on the third-generation mirror quark mass $M_3$. We can see that the branching ratio of $t \to cgH$ increases as $M_3$ increases; the largest branching ratio comes from case III, and the maximum value can reach $1.4 \times 10^{-5}$. In Fig. 6(b), we show the dependence of the branching ratio of the $t \to cgH$ decay process on the scale $f$. We can see that the correction of the LHT model decouples as the scale $f$ increases.

In Fig. 7, we show the dependence of the branching ratio of the $t \to c\gamma H$ decay process on the third-generation mirror quark mass $M_3$ (a) and the scale $f$ (b). We can see that the decay process $t \to c\gamma H$ has similar behaviors as the decay process $t \to cgH$. The maximum value of the branching ratio can reach $3.5 \times 10^{-7}$.

## V. SUMMARY

In this paper, we calculated the top quark FCNC decays $t \to cH$ and $t \to cg(\gamma)H$ in the unitary gauge in the LHT model. We found that the branching ratio for $t \to cH$ and $t \to cg(\gamma)H$ can respectively reach $5.8 \times 10^{-5}$ and $1.4 \times 10^{-5}$ $(3.5 \times 10^{-7})$ in the allowed parameter space. When the mirror fermion mass $M_3 > 2.2$ TeV and the cutoff scale $f = 500$ GeV, $t \to cH$ can reach $3\sigma$ sensitivity at 8 TeV LHC with luminosity $\mathcal{L} = 20\ \mathrm{fb}^{-1}$. We also noted that the 14 TeV LHC has the potential to observe this channel at the $5\sigma$ sensitivity level for $M_3 = 2.1(1.5)$ TeV when $\mathcal{L} = 30(300)\ \mathrm{fb}^{-1}$. Therefore, we can see that $t \to cH$ may be used to test the LHT model at the LHC.

## ACKNOWLEDGEMENT

We appreciate helpful discussions with Jinmin Yang, Junjie Cao, Lei Wang and Lei Wu. N.L. would like to thank Dr. Archil Kobakhidze for his warm hospitality in the Sydney node of CoEPP in Australia. This work is supported by the National Natural Science Foundation of China (NNSFC) under Grants No. 11275057, No. 11305049, and No. 11347140, and the Specialized Research Fund for the Doctoral Program of Higher Education under Grant No. 20134104120002.

[1] S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D **2**, 1285 (1970).

[2] G. Eilam, J. L.Hewett, and A. Soni, Phys. Rev. D **59**, 039901(E) (1998); T. P. Cheng and M. Sher, Phys. Rev. D **35**, 3484 (1987); B. Grzadkowski, J. F. Gunion, and P. Krawczyk, Phys. Lett. B **268**, 106 (1991); A. Antaramian, L. J. Hall, and A. Rasin, Phys. Rev. Lett. **69**, 1871 (1992); N. G. Deshpande, B. Margolis, and H. D. Trottier, Phys. Rev. D **45**, 178 (1992); W.-S. Hou, Phys. Lett. B **296**, 179 (1992); M. Luke and M. J. Savage, Phys. Lett. B **307**, 387 (1993); L. Hall and S. Weinberg, Phys. Rev. D **48**, R979 (1993); D. Atwood, L. Reina, and A. Soni, Phys. Rev. Lett. **75**, 3800 (1995); , Phys. Rev. D **53**, 1199 (1996); , Phys. Rev. D **55**, 3156 (1997).

[3] J. A. Aguilar-Saavedra, Nucl. Phys. **B843**, 638 (2011) [Nucl. Phys.**B851**, 443(E) (2011)]; J. A. Aguilar-Saavedra, Acta Phys. Pol. B **35**, 2671 (2004); Phys. Lett. B **502**, 115 (2001); J. A. Aguilar-Saavedra and G. C. Branco, Phys. Lett. B **495**, 347 (2000); F. del Aguila and J. A. Aguilar-Saavedra, Nucl. Phys. **B576**, 56 (2000); F. del Aguila, J. A. Aguilar-Saavedra, and L. Ametller, Phys. Lett. B **462**, 310 (1999).

[4] J. Drobnak, S. Fajfer, and J. F. Kamenik, Phys. Rev. D **82**, 073016 (2010); , Phys. Rev. Lett. **104**, 252001 (2010).

[5] P. M. Ferreira, R. B. Guedes, and R. Santos, Phys. Rev. D **77**, 114008 (2008); R. A. Coimbra, P. M. Ferreira, R. B. Guedes, O. Oliveira, A. Onofre, R. Santos, and M. Won, Phys. Rev. D **79**, 014006 (2009); P. M. Ferreira and R. Santos, Phys. Rev. D **74**, 014006 (2006); P. M. Ferreira and R. Santos, *ibid.* **73**, 054025 (2006); P. M. Ferreira, O. Oliveira, and R. Santos, Phys. Rev. D **73**, 034011 (2006).

[6] C. Han, N. Liu, L. Wu, and J. M. Yang, Phys. Lett. B **714**, 295 (2012); J. Cao, L. Wu, and J. M. Yang, Phys. Rev. D **83**, 034024 (2011); J. Cao, L. Wang, L. Wu, and J. M. Yang, Phys. Rev. D **84**, 074001 (2011); N. Liu and L. Wu, Commun. Theor. Phys. **55**, 296 (2011); J. Cao, Z. Heng, L. Wu, and J. M. Yang, Phys. Rev. D **81**, 014016 (2010); G.-R. Lu and L. Wu, Chin. Phys. Lett. **27**, 031401 (2010).

[7] C. S. Li, R. J. Oakes, and J. M. Yang, Phys. Rev. D **49**, 293 (1994); , Phys. Rev. D **56**, 3156 (1997); J. L. Lopez, D. V. Nanopoulos, and R. Rangarajan, Phys. Rev. D **56**, 3100 (1997); G. M. de Divitiis, R. Petronzio, and L. Silvestrini, Nucl. Phys. **B504**, 45 (1997); J. Guasch and J. Sola, Nucl. Phys. **B562**, 3 (1999); J. J. Liu, C. S. Li, L. L. Yang, and L. G. Jin, Phys. Lett. B **599**, 92 (2004).

[8] G. Eilam, A. Gemintern, T. Han, J. M. Yang, and X. Zhang, Phys. Lett. B **510**, 227 (2001).

[9] J. Drobnak, S. Fajfer, and J. F. Kamenik, J. High Energy Phys. 03 (**2009**) 077; A. Arhrib, Phys. Rev. D **72**, 075016 (2005); S. Bejar, J. Guasch, and J. Sola, Nucl. Phys. **B675**, 270 (2003); E. O. Iltan, Phys. Rev. D **65**, 075017 (2002); W. S. Hou, Phys. Lett. B **296**, 179 (1992); R. A. Daz, R. Martnez, and J. Alexis Rodrguez, arXiv:hep-ph/0103307.

[10] C. Zhang and F. Maltoni, Phys. Rev. D **88**, 054005 (2013).

[11] S. Bar-Shalom, G. Eilam, M. Frank, and I. Turan, Phys. Rev. D **72**, 055018 (2005); J. L. Díaz-Cruz, M. A. Pérez, G. Tavares-Velasco, and J. J. Toscano, Phys. Rev. D **60**, 115014 (1999); E. O. Iltan and I. Turan, Phys. Rev. D **67**, 015004 (2003); S. Bar-Shalom, G. Eilam, A. Soni, and J. Wudka, Phys. Rev. D **57**, 2957 (1998).

[12] G. Eilam, M. Frank, and I. Turan, Phys. Rev. D **74**, 035012 (2006); J. J. Cao, G. Eilam, M. Frank, K. Hikasa, G. L. Liu, I. Turan, and J. M. Yang, Phys. Rev. D **75**, 075021 (2007); Z. Heng, G. Lu, L. Wu, and J. M. Yang, Phys. Rev. D **79**, 094029 (2009).

[13] J. Han, B. Li, and X. Wang, Phys. Rev. D **83**, 034032 (2011).

[14] ATLAS Collaboration, Report No. ATLAS-CONF-2013-081.

[15] H. C. Cheng and I. Low, J. High Energy Phys. 09 (**2003**) 051; J. High Energy Phys. 08 (**2004**) 061; I. Low, J. High Energy Phys. 10 (**2004**) 067; J. Hubisz and P. Meade, Phys. Rev. D **71**, 035016 (2005).

[16] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson, J. High Energy Phys. 07 (**2002**) 034; S. Chang, J. High Energy Phys. 12 (**2003**) 057; T. Han, H. E. Logan, B. McElrath, and L. T. Wang, Phys. Rev. D **67**, 095004 (2003); M. Schmaltz and D. Tucker-smith, Annu. Rev. Nucl. Part. Sci. **55**, 229 (2005).

[17] X. F. Han, L. Wang, and J. M.Yang, Phys. Rev. D **80**, 015018 (2009); J. Han, B. Li, and X. Wang, Phys. Rev. D **83**, 034032 (2011).

[18] J. Hubisz, S. J. Lee, and G. Paz, J. High Energy Phys. 06 (**2006**) 041; M. Blanke A. J. Buras, S. Recksiegel, C. Tarantino, arXiv:0805.4393; M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig, and A. Weiler, J. High Energy Phys. 01 (**2007**) 066; M. Blanke, A. J. Buras, A. Poschenrieder, C. Tarantino, S. Uhlig, and A. Weiler, J. High Energy Phys. 12 (**2006**) 003.

[19] X. F. Han, L. Wang, and J. M. Yang, Phys. Rev. D **78**, 075017 (2008); C. X. Yue, J. Y. Liu, and S. H. Zhu, Phys. Rev. D **78**, 095006 (2008).

[20] G. Aad *et al.* (ATLAS Collaboration), Phys. Rev. Lett. **108**, 041805 (2012).

[21] J. A. Aguilar-Saavedra, Acta Phys. Pol. B **35**, 2671 (2004).

[22] M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig, and A. Weiler, Phys. Lett. B **646**, 253 (2007).

[23] Y.-J. Zhou, H.-S. Hou, and H. Sun, Commun. Theor. Phys. **59**, 443 (2013).

[24] M. Blanke, A. J. Buras, A. Poschenrieder, C. Tarantino, S. Uhlig, and A. Weiler, J. High Energy Phys. 12 (**2006**) 003; M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig, and A. Weiler, J. High Energy Phys. 01 (**2007**) 066.

[25] F. del Aguila, M. Pérez-Victoria, and J. Santiago, J. High Energy Phys. 09 (**2000**) 011.

[26] G. Passarino and M. J. G. Veltman, Nucl. Phys. **B160**, 151 (1979).

[27] G. 't Hooft and M. J. G. Veltman, Nucl. Phys. **B153**, 365 (1979).

[28] T. Hahn and M. Perez-Victoria, Comput. Phys. Commun. **118**, 153 (1999); T. Hahn, Nucl. Phys. B, Proc. Suppl. **135**, 333 (2004).

[29] T. Goto, Y. Okada, and Y. Yamamoto, Phys. Lett. B **670**, 378 (2009).

[30] M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig, and A. Weiler, J. High Energy Phys. 01 (**2007**) 066; J. High Energy Phys. 12 (**2006**) 003.

[31] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).

[32] B. Yang, X. Wang, and J. Han, Nucl. Phys. **B847**, 1 (2011); J. Hubisz, P. Meade, A. Noble, and M. Perelstein, J. High Energy Phys. 01 (**2006**) 135; A. Belyaev, C.-R. Chen, K. Tobe, and C. P. Yuan, Phys. Rev. D **74**, 115020 (2006); Q.-H. Cao and C.-R. Chen, Phys. Rev. D **76**, 075007 (2007); L. Fiorini, arXiv:1201.5844; K. Blum, Y. Grossman, Y. Nir, and G. Perez, Phys. Rev. Lett. **102**, 211802 (2009); J. Reuter and M. Tonini, J. High Energy Phys. 13 (**2013**) 077; X.-F. Han, L. Wang, J. M. Yang, and J. Zhu, Phys. Rev. D **87**, 055004 (2013); J. Reuter, M. Tonini, and M. de Vries, arXiv:1307.5010.

[33] J. Hubisz, P. Meade, A. Noble, and M. Perelstein, J. High Energy Phys. 01 (**2006**) 135; M. Perelstein and J. Shao, Phys. Lett. B **704**, 510 (2011); G. Cacciapaglia, A. Deandrea, J. Ellis, J. Marrouche, and L. Panizzi, Phys. Rev. D **87**, 075006 (2013).

[34] B. Mele, S. Petrarca, and A. Soddu, Phys. Lett. B **435**, 401 (1998).

[35] J. A. Aguilar-Saavedra, Acta Phys. Pol. B **35**, 2671 (2004).

[36] D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **55**, 3156 (1997); S. Bejar, J. Guasch, and J. Sola, Nucl. Phys. **B600**, 21 (2001).

[37] J. Guasch and J. Sola, Nucl. Phys. **B562**, 3 (1999).

[38] J. Cao, G. Eilam, M. Frank, K. Hikasa, G. Liu, I. Turan, and J. Yang, Phys. Rev. D **75**, 075021 (2007).

[39] C. Kao, H.-Y. Cheng, W.-S. Hou, and J. Sayre, Phys. Lett. B **716**, 225 (2012); K.-F. Chen, W.-S. Hou, C. Kao, and M. Kohda, Phys. Lett. B **725**, 378 (2013).

[40] P. P. Giardino, K. Kannike, I. Masina, M. Raidal, and A. Strumia, arXiv:1303.3570.