Flavor changing top quark decay and bottom-strange production in the littlest Higgs model with T-parity

Zhou Ya-Jin\textsuperscript{a}\textsuperscript{*}, Hou Hong-Sheng\textsuperscript{b}\textsuperscript{†}, Sun Hao\textsuperscript{c}\textsuperscript{‡}

\textsuperscript{a} School of Physics, Shandong University, Jinan Shandong 250100, P.R. China
\textsuperscript{b} Department of Physics, Hangzhou Normal University, Hangzhou Zhejiang 310036, P.R.China
\textsuperscript{c} School of Physics and Technology, University of Jinan, Jinan Shandong 250022, P.R.China

Abstract

Flavor changing effects on the processes $t \rightarrow cH$, $e^{+}e^{-} \rightarrow b\bar{s}$, $e^{+}e^{-} \rightarrow b\bar{s}H$ and $pp \rightarrow b\bar{s}$ in the LHT model are investigated in this paper. We calculate the one-loop level contributions from the T-parity odd mirror quarks and gauge bosons. The results show that the top quark rare decay $t \rightarrow cH$ in the LHT model can be significantly enhanced relative to that in the SM. The $b\bar{s}$ production at linear colliders in the LHT model can enhance the SM cross section a lot and reach 0.1 fb in some parameter space allowed in the experiment. But the heavy gauge boson and mirror fermion loops have small contribution to the processes $pp \rightarrow b\bar{s}$ and $e^{+}e^{-} \rightarrow b\bar{s}H$. So the LHT effect on $e^{+}e^{-} \rightarrow b\bar{s}$ might be detected at future linear colliders, while it’s too small to be seen for the $e^{+}e^{-} \rightarrow b\bar{s}H$ and $pp \rightarrow b\bar{s}$ processes at future linear colliders and LHC.

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I Introduction

The Flavor Changing Neutral Currents (FCNC) couplings play an important role in searching for new physics beyond the Standard Model (SM) for the following reasons: they are forbidden at tree level and suppressed by the GIM mechanism [1] at one loop level in the SM, on the other hand, many new particles appear in the loops in new physics models beyond SM, which may enhance the flavor changing transitions [2]. Furthermore searching for FCNC is one of the main goals of the high energy colliders.

The Little Higgs mechanism [3,4] offers a solution to the hierarchy problem without fine tuning. The most compact implementation of the Little Higgs mechanism is known as the Littlest Higgs model (LH) [5,6]. But the original Little Higgs models suffer strong constraints from electroweak precision data [7]. To solve this problem, a $\mathbb{Z}_2$ discrete symmetry named “T-Parity” is introduced in Ref. [8,9], by which dangerous diagrams with the tree level exchange of heavy neutral gauge bosons are forbidden. The Littlest Higgs model with T-parity (LHT) requires the introduction of “mirror fermions” for each SM fermion doublet. The mirror fermions are T-odd and can obtain large masses. These mirror fermions and the heavy gauge bosons appear in loop diagrams can induce new contributions to quark FCNC processes.

FCNC effects are usually studied via decay mode, such as rare $B, K, D$ meson decay, top quark rare decay, Higgs and $Z$ decay, etc. Besides via decay modes, the flavor changing vertices $bsV (V = \gamma, Z)$ can also be investigated via bottom-strange associated production. The $e^+e^- \rightarrow b\bar{s}$ process in the SM have been studied in [10], and the results show that the cross section is larger than that of the $t\bar{c}$ production. We checked these two processes in the LHT model, and found that this character is kept. We will study the flavor changing effects via $b\bar{s}$ and $b\bar{s}H$ productions at linear colliders in the LHT model. Since the phase space of $b\bar{s}$ production at the large hadron collider (LHC) is very large, we will investigate this process as well and looking forward to get large contribution. On the other side, the top quark plays a special role in FCNC phenomenology due to its heaviness. The decay mode $t \rightarrow cH$ is forbidden at tree level, and its decay width is about $10^{-14}$ at one loop
level in the SM \cite{11}. In the MSSM the width could be enhanced to $10^{-5} \sim 10^{-4}$ \cite{12}. In Ref. \cite{13} the $t \rightarrow cH$ process is studied in the LH model, and the authors found that the branch ratio is at most of the order $10^{-12}$. With more mirror particles in the loops and less constraints on the parameters, $t \rightarrow cH$ process in the LHT model would have much larger decay width than in the LH model.

In this paper we investigate the flavor changing effects on the following processes: $t \rightarrow cH$, $e^+e^- \rightarrow b\bar{s}$, $e^+e^- \rightarrow b\bar{s}H$ and $pp \rightarrow b\bar{s}$. Since the final states $b\bar{s}$ and $b\bar{s}$ are undistinguishable in the experiment, we need to calculate the two individual processes $b\bar{s}(H)$ and $b\bar{s}(H)$ and sum their cross sections up. We found that the two individual processes have the same cross sections, so we only give the results for $b\bar{s}(H)$ final state, the total cross sections can be obtained by doubling the $b\bar{s}(H)$ ones. The paper is organized as follows: In Sec.II we have a brief review of the LHT model. In Sec.III we give the analytical and numerical calculations. Finally a short summary is given.

II Brief review of the model

Here we briefly review the structure of the LHT model relevant to our analysis, and the detailed description can be found in the literature \cite{14, 15}.

II.1 Gauge and Higgs sectors

Gauge and Higgs sectors of the littlest Higgs model are described as a nonlinear $\sigma$ model with the spontaneous global symmetry breaking from SU(5) to SO(5) with scalar fields. From the SU(5)/SO(5) breaking, there arise 14 Goldstone bosons which are described by the “pion” matrix $\Pi$, given explicitly by

$$
\Pi = \begin{pmatrix}
-\frac{\omega^0}{2} - \frac{\eta}{\sqrt{20}} & -\frac{\omega^+}{\sqrt{2}} & -i\phi^+ & -i\phi^+ & -i\phi^+ \\
-\omega^+ & -\frac{\eta}{2} & v+h+i\eta^0 & 0 & 0 \\
\frac{\pi}{\sqrt{2}} & \frac{\pi}{\sqrt{2}} & -\frac{\omega^+}{\sqrt{2}} & \frac{v+h+i\eta^0}{2} & \frac{v+h+i\eta^0}{2} \\
\frac{\phi^-}{\sqrt{2}} & \frac{\phi^-}{\sqrt{2}} & \frac{\phi^-}{\sqrt{2}} & \frac{\phi^-}{\sqrt{2}} & \frac{\phi^-}{\sqrt{2}} \\
\frac{\phi^0}{\sqrt{2}} & \frac{\phi^0}{\sqrt{2}} & \frac{\phi^0}{\sqrt{2}} & \frac{\phi^0}{\sqrt{2}} & \frac{\phi^0}{\sqrt{2}} \\
\end{pmatrix}
$$

(2.1)
where it consists of a doublet $H$ and a triplet $\Phi$ under the unbroken $SU(2)_L \times U(1)_Y$ group which are given by

$$H = \left( \frac{-i \pi^+}{\sqrt{2}(v + h + i\pi^0)} \right), \quad \Phi = \left( \begin{array}{c} -i\phi^+ \frac{v}{\sqrt{2}} + i\phi^0 \frac{h}{\sqrt{2}} \\ -i\phi^0 \frac{v}{\sqrt{2}} - i\phi^+ \frac{h}{\sqrt{2}} \end{array} \right).$$

(2.2)

Here, $H$ plays the role of the SM Higgs doublet, $h$ is the physical Higgs field and $v \simeq 246$ GeV. The fields $\eta$ and $\omega$ are eaten by new heavy gauge bosons $A_H, W_H$ and $Z_H$ when the $[SU(2) \times U(1)]^2$ gauge group is broken down to $SU(2)_L \times U(1)_Y$, whereas the $\pi$ fields are absorbed by the standard model $W/Z$ bosons after electroweak symmetry breaking (EWSB). The field $h$ and $\Phi$ remain in the spectrum.

Under the T-parity, the SM particles are T-even, and all the new particles except $T_+$ (as we will introduce later) are T-odd. The masses of T-even gauge bosons are given to $v^2/f^2$ order by:

$$M_{W_L} = \frac{gv}{2}(1 - \frac{v^2}{12f^2}), \quad M_{Z_L} = \frac{MW_L}{\cos \theta_W}, \quad MA_L = 0,$$

(2.3)

where $\theta_W$ is the weak mixing angle, and $g$ is the SM $SU(2)$ gauge couplings.

The masses of T-odd gauge bosons are given to $v^2/f^2$ order by:

$$M_{W_H} = g f (1 - \frac{v^2}{8f^2}), \quad M_{Z_H} = M_{W_H}, \quad MA_H = \frac{g'f}{\sqrt{5}}(1 - \frac{5v^2}{8f^2})$$

(2.4)

where $g'$ is the SM U(1) gauge couplings.

II.2 Fermion sector

The T-even fermion sector includes the SM quarks, leptons, and an additional heavy quark $T_+$, which is introduced in the LH model in order to cancel the quadratic divergence of the Higgs mass coming from top loops. Each T-even fermion need a T-odd mirror fermion under T-parity. The mirror quarks and leptons are involved in this paper. We denote them by

$$\left( \begin{array}{c} u^1_H \\ d^1_H \end{array} \right), \quad \left( \begin{array}{c} u^2_H \\ d^2_H \end{array} \right), \quad \left( \begin{array}{c} u^3_H \\ d^3_H \end{array} \right),$$

(2.5)

and

$$\left( \begin{array}{c} \nu^1_H \\ l^1_H \end{array} \right), \quad \left( \begin{array}{c} \nu^2_H \\ l^2_H \end{array} \right), \quad \left( \begin{array}{c} \nu^3_H \\ l^3_H \end{array} \right).$$

(2.6)
In $\mathcal{O}(v^2/f^2)$ their masses are given by
\begin{align}
m_{u_i^H} &= \sqrt{2}\kappa_q^i f (1 - \frac{v^2}{8f^2}) = m_{H_i}(1 - \frac{v^2}{8f^2}), \\
m_{d_i^H} &= \sqrt{2}\kappa_q^i f = m_{H_i} \\
m_{\nu^i_H} &= \sqrt{2}\kappa_l^i f (1 - \frac{v^2}{8f^2}) = m_{H_i}(1 - \frac{v^2}{8f^2}), \\
m_{l^i_H} &= \sqrt{2}\kappa_l^i f = m_{H_i}
\end{align}
where $i$ is the generation index, and $\kappa_q^i$ ($\kappa_l^i$) are the eigenvalues of the mirror quark (lepton) Yukawa coupling matrices. We neglect $\mathcal{O}(v^2/f^2)$ differences between $m_{u_i^H}$ and $m_{d_i^H}$ ($m_{\nu^i_H}$ and $m_{l^i_H}$) in the numerical calculation because these differences only contribute to higher order corrections in the $v^2/f^2$ expansion.

II.3 T-odd flavor mixing

In the LHT model, the mirror fermions open up a new flavor structure in the model. As discussed in Ref. [15, 17], there are four CKM-like unitary mixing matrices in the mirror fermion sector: $V_{Hu}$, $V_{Hd}$, $V_{Hl}$ and $V_{H\nu}$. These mirror mixing matrices are involved in the flavor changing interactions between SM fermions and T-odd mirror fermions which are mediated by the T-odd heavy gauge and Goldstone bosons ($W_H, Z_H, A_H$ and $\omega^\pm, \omega^0, \eta$), and they satisfy
\begin{align}
V_{Hu}^\dagger V_{Hd} &= V_{\text{CKM}}, \\
V_{H\nu}^\dagger V_{Hl} &= V_{\text{PMNS}}. 
\end{align}

Using the method in Ref. [17, 18], $V_{Hd}$ is parameterized with three angles $\theta_{12}^d, \theta_{23}^d, \theta_{13}^d$ and three phases $\delta_{12}^d, \delta_{23}^d, \delta_{13}^d$, and analogously $V_{Hl}$ is parameterized with three angles $\theta_{12}^l, \theta_{23}^l, \theta_{13}^l$ and three phases $\delta_{12}^l, \delta_{23}^l, \delta_{13}^l$. The explicit expression won’t be listed here. The Feynman rules for the flavor violating interactions which are involved in our analysis can be found in Ref. [17, 19].

III Calculation

In this section we calculate the flavor changing effects originated from heavy gauge bosons and mirror quarks on the processes $t \rightarrow cH$, $e^+e^- \rightarrow b\bar{s}$, $e^+e^- \rightarrow b\bar{s}H$ and $pp \rightarrow b\bar{s}$ in the LHT model. The relevant Feynman diagrams are shown in Figs[18]. The diagrams
for the subprocesses of $pp \rightarrow b\bar{s}$ are similar to the diagrams of process $e^+e^- \rightarrow b\bar{s}$ but much more, so we don’t list them in the paper. We have added the relevant Feynman rules of the LHT model to _FeynArts3_ package [20] and use it to generate the Feynman diagrams and the corresponding amplitudes. In the calculations of the one-loop diagrams we adopt the definitions of one-loop integral functions as in Ref. [21]. The loop integral functions are calculated by using the formulas in Ref. [22].

![Feynman Diagrams](image)

Figure 1: Feynman diagrams for $t \rightarrow cH$ in the LHT model.

Since there are no tree level diagrams for these processes except for $pp \rightarrow u\bar{u} \rightarrow b\bar{s}$, we just sum all the unrenormalized reducible and irreducible one-loop diagrams, and the results will be finite and gauge-invariant. We checked these processes ($pp \rightarrow u\bar{u} \rightarrow b\bar{s}$ not included) and found that the divergences are canceled at $\mathcal{O}(v^2/f^2)$ for all the processes except $t \rightarrow cH$. This divergence was explained as the sensitivity of the decay amplitudes to the UV completion of the LH model in Ref. [17, 23], and it is gauge independent. We use the ’t Hooft gauge in our calculation. After update some vertices to $\mathcal{O}(v^2/f^2)$ in Ref. [19], Goto et al. and Blanke et al. found that the logarithmic divergence in Z boson flavor changing processes can be canceled. We calculated the process $t \rightarrow cH$ with the updated Feynman rules, and found that the divergence can’t be canceled. So in our numerical calculations, we remove the divergent term $1/\epsilon$ and take the renormalization scale $\mu = \Lambda$ with $\Lambda = 4\pi f$ being the cutoff scale of the LHT model, as in Ref. [17, 23].

In the numerical calculation, we take the SM parameters as follows [24–26]

$$\alpha = 1/128, \quad \alpha_s(m_Z) = 0.1184, \quad m_W = 80.385 \text{ GeV}, \quad m_Z = 91.1876 \text{ GeV},$$

$$\Gamma_Z = 2.4952 \text{ GeV}, \quad m_t = 173.5 \text{ GeV}, \quad m_c = 1.27 \text{ GeV}, \quad m_b = 4.67 \text{ GeV},$$

$$m_s = 0.101 \text{ GeV}, \quad m_H = 125 \text{ GeV}. \quad (3.1)$$
Figure 2: Flavor changing vertex $\gamma(Z)b\bar{s}$ in the LHT model.

For the scattering processes, we take the cuts on final particles as $p_T^{b,s,H} \geq 15$ GeV.

Considering the constraints on PMNS matrix [27–30], we set PMNS parameters to

$$s_{12} = \sqrt{0.3}, \ s_{23} = \sqrt{0.5}, \ s_{13} = \sqrt{0.024}, \ \delta = 65^\circ,$$

where the Majorana phases in $V_{PMNS}$ have been set to zero, because no Majorana mass term has been introduced for right-handed neutrinos.

The LHT parameters which are relevant to our analysis are

$$f, \ m_{H1}, \ m_{H2}, \ m_{H3}, \ \theta^d_{12}, \ \theta^d_{23}, \ \theta^d_{13}, \ \delta^d_{12}, \ \delta^d_{23}, \ \delta^d_{13}$$

$$m^l_{H1}, \ m^l_{H2}, \ m^l_{H3}, \ \theta^l_{12}, \ \theta^l_{23}, \ \theta^l_{13}, \ \delta^l_{12}, \ \delta^l_{23}, \ \delta^l_{13}$$

The LHT scale $f$ can be as low as 500 GeV [31], so we vary it in the range $500 \text{ GeV} \leq f \leq 1500$ GeV. The constrains on the mass spectrum of the mirror fermions have been extensively studied [15–17, 19, 32–35]. It is convenient to consider several representative scenarios for the structure of the matrix $V_{Hd}$. In Ref [16,17] several benchmark scenarios was introduced, among which Scenario 4 allows for large effects in the $B_s$ system. In this scenario the hierarchical structure of $V_{Hd}$ matrix is very different from the structure of
Figure 3: One-loop Feynman diagrams for process $e^+e^- \rightarrow b\bar{s}H$ and $e^+e^- \rightarrow b\bar{s}$ in the LHT model. The loop-induced $\gamma(Z)b\bar{s}$ vertex (gray circle in (a) and (b)) is shown in Fig.2.

the CKM matrix, and they assume that

$$\frac{1}{\sqrt{2}} \leq s_{12} \leq 0.99, \ 5 \times 10^{-5} \leq s_{23}^d \leq 2 \times 10^{-4}, \ 4 \times 10^{-2} \leq s_{13}^d \leq 0.6. \quad (3.5)$$

To be simplicity we choose the lower and upper limits of $s_{ij}^d$ ($(ij) = (12), (23), (13)$) in Eq.(3.5) as Case II and Case III, respectively, with the phase term $\delta_{ij}^d = 0$. In Case I we assume that there are no mixing in down type mirror quarks, i.e., $V_{Hd} = 1$. We follow Ref [36] to give the values of mirror fermions. Here we list the there cases we used,

- **Case I**, $V_{Hd} = 1$

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

- **Case III**, $s_{12}^d = 0.99, \ s_{23}^d = 2 \times 10^{-4}, \ s_{13}^d = 0.6, \ \delta_{12}^d = \delta_{23}^d = \delta_{13}^d = 0$
Figure 4: Branching ratios for \( t \rightarrow cH \) in Case I, as the function of \( m_{H3} \) (left) and \( f \) (right).

In all the three cases the first two mirror quark generations are chosen to be quasi-generate. In Case II and III we take the masses of mirror fermions as

\[
m_{d1} = m_{d2} = \frac{600 \text{ GeV}}{\text{TeV}} f, \quad m_{d3} = \frac{1400 \text{ GeV}}{\text{TeV}} f.
\] (3.6)

The lepton sector in Eq.(3.4) is only involved in the process \( e^+e^- \rightarrow b\bar{s} \). According to Scenario C in Ref. [37], we constrain the mirror lepton masses to lie in the range \( 300 \text{ GeV} \leq m_{lH} \leq 1.5 \text{ TeV} \), and scan over the whole parameter space for the mirror lepton mixing part, i.e. \( 0 \leq \theta_{ij} \leq 2\pi \) and \( 0 \leq \delta_{ij} \leq 2\pi \).

For the decay process \( t \rightarrow cH \), the branching ratio is defined as follow because \( t \rightarrow bW^+ \) is the dominant channel of top-quark

\[
Br(t \rightarrow cH) = \frac{\Gamma(t \rightarrow cH)}{\Gamma(t \rightarrow bW^+)}.
\] (3.7)

For the other processes we show the cross sections as final results.

Case I only contribute to \( t \rightarrow cH \) and \( pp \rightarrow u\bar{u} \rightarrow b\bar{s} \) processes, however the cross section for the latter is so small (\( \sim 10^{-7} \) fb) that can be neglected. The other processes have no contribution from mirror quark loops, because there are no flavor mixing between down-type mirror quarks in Case I.

In Fig.4 we show the branching ratios of \( t \rightarrow cH \) decay process as the functions of the mass of the third generation mirror quark (left) and the LHT scale \( f \) (right) in
Figure 5: Branching ratios for $t \rightarrow cH$ as the function of $f$ in Case II and III. Case I. We set $m_{H_1} = m_{H_2} = 300$ GeV and $f = 500$ GeV in the left figure. The branching ratio increases with $m_{H_3}$, because the decay rate is enhanced by the mass splitting between the three generation mirror quarks, and since we set $m_{H_1} = m_{H_2}$, there is only one mass splitting $m_{H_3} - m_{H_2}$, which increases with $m_{H_3}$. In the right figure we set $m_{H_1} = m_{H_2} = 0.6f$ and $m_{H_3} = 3f$, and we can see that the branching ratio decreases with $f$, but very slowly. That’s because the mass splitting $m_{H_3} - m_{H_2}$ is large enough ($2.4f$) to cancel the the decrement caused by large $f$. We can also see that the branching ratio enhances a lot in the LHT model compared with that in the SM ($10^{-14}$). At the same time the contributions from the LHT model are much larger than those in the LH model without T-parity ($\sim 10^{-12}$). That’s because the parameters have less constraints, and the mirror fermions in the loops contribute a lot. The branching ratio can reach $10^{-4}$ when $f=500$ GeV, $m_{H_1} = m_{H_2} = 300$ GeV and $m_{H_3} = 3$ TeV, which is even larger than the SUSY-QCD contribution.

In Fig.5 we present the branching ratio of $t \rightarrow cH$ process as the function of $f$ in Case II and III. In both cases the branching ratios decrease with $f$ but very slowly, for the same reason stated above. The branching ratio in Case III is much larger than that in Case II, because the mixing between mirror quarks is much larger.

In Fig.6 we present the cross sections for processes $e^+e^- \rightarrow b\bar{s}$, $e^+e^- \rightarrow b\bar{s}H$ and $pp \rightarrow b\bar{s}$ as the functions of $f$ for Case II (left) and III (right). The center of mass system (c.m.s.) energy $\sqrt{s}$ is 14 TeV at LHC and 500 GeV at a linear collider. We scanned
the mirror lepton sector parameters over the whole space except the restriction for mirror lepton masses when computing the $e^+e^- \rightarrow b\bar{s}$ cross sections. In Case II the cross sections for $e^+e^- \rightarrow b\bar{s}$ process in the LHT model almost degenerate with those in the SM in most of the parameter space, only a small enhancement in the low $f$ region. While in Case III the LHT effect could enhance the SM cross section by two orders, with $f$ close to 500 GeV. In both cases the LHT effects on $e^+e^- \rightarrow b\bar{s}H$ process are very small, only enhance the SM cross section by a few percent.

Coming to $pp \rightarrow b\bar{s}$ process, there is a tree-level diagram of flavor changing charged current for the subprocess $u\bar{u} \rightarrow b\bar{s}$. If we compute one-loop contribution, there would exist UV and IR divergences, so the renormalization procedure is necessary. For simplicity we estimate the cross section in the SM, and compute the pure LHT effect on this process at 1-loop level. This means that we set the CKM matrix to be unit when we generate the Feynman diagrams and amplitudes in the LHT model, and in this case there are no tree level diagrams, so we can sum all the diagrams together and the result would be finite. Now let’s make a rough comparison between the cross section in the SM and in the LHT model. First we list the cross sections for the three subprocesses at LHC:

\[
\sigma^{tree}_{SM}(pp \rightarrow u\bar{u} \rightarrow b\bar{s}) = 0.449 \text{ fb} \quad (3.8)
\]
\[
\sigma^{1\text{-}loop}_{SM}(pp \rightarrow d\bar{d} \rightarrow b\bar{s}) = 0.505 \text{ fb} \quad (3.9)
\]
\[ \sigma_{SM}^{1-loop}(pp \rightarrow gg \rightarrow b\bar{s}) = 0.604 \text{ fb} \] (3.10)

Supposing the QCD correction of \( pp \rightarrow u\bar{u} \rightarrow b\bar{s} \) to be 20%, we got

\[ \sigma_{SM}^{1-loop}(pp \rightarrow u\bar{u} \rightarrow b\bar{s}) \sim 0.449 \times 0.2 \text{ fb} = 0.898 \text{ fb} \] (3.11)

Summing the tree level and 1-loop level subprocess cross sections together, we obtain the total cross section of \( pp \rightarrow b\bar{s} \) process in the SM at 1-loop level

\[ \sigma_{SM}^{1-loop}(pp \rightarrow b\bar{s}) \sim 1.65 \text{ fb} \] (3.12)

From Fig. 6(a) we can see that the pure LHT cross section of \( pp \rightarrow b\bar{s} \) varies from \( 3 \times 10^{-5} \) to \( 3 \times 10^{-3} \) fb, which is much smaller than the SM contribution. In Fig. 6(b) the pure LHT cross section for \( pp \rightarrow b\bar{s} \) enhances 2 orders compare with Fig. 6(a), but still smaller than the SM cross section, and can’t be detected at LHC.

In Fig. 7 we present the cross sections for \( e^+e^- \rightarrow b\bar{s} \) and \( e^+e^- \rightarrow b\bar{s}H \) as the functions of \( \sqrt{s} \) in Case III, with \( f = 500 \) GeV, \( m_{H1} = m_{H2} = 300 \) GeV and \( m_{H3} = 700 \) GeV. There are three peaks in the curves for \( e^+e^- \rightarrow b\bar{s} \) process, corresponding to the resonance of Z boson, a pair of W boson threshold, and a pair of top quark threshold, respectively. The cross section for \( e^+e^- \rightarrow b\bar{s}H \) decreases with the increase of \( \sqrt{s} \) beyond the Higgs resonance peak (250 GeV) because of the s-channel depression. We can also see that the LHT effect on \( e^+e^- \rightarrow b\bar{s} \) process increase with \( \sqrt{s} \), and could enhance the SM cross section by 3 orders when \( \sqrt{s} \gtrsim 800 \text{GeV} \). The cross section can reach 0.1 fb with large \( \sqrt{s} \), and even exceed 1 fb at Z resonance. So it might be possible to see the LHT effect on this process at future linear colliders. While the cross section in the LHT model for \( e^+e^- \rightarrow b\bar{s}H \) process almost degenerates with that in the SM, and they are too small to be detected at the future linear colliders.

**IV Summary**

In this paper we calculate the one-loop contributions from heavy gauge bosons and mirror fermions to the top quark rare decay process \( t \rightarrow cH, b\bar{s}(H) \) production at linear colliders,
Figure 7: Cross sections for $e^+e^- \to b\bar{s}$ and $e^+e^- \to b\bar{s}H$ as the functions of $\sqrt{s}$ in Case III, with $f = 500$ GeV, $m_{H_1} = m_{H_2} = 300$ GeV and $m_{H_3} = 700$ GeV.

and $b\bar{s}$ production at the 14 TeV LHC in the LHT model. The branching ratio for $t \to cH$ can reach $10^{-4}$ in the parameter space we considered, which is much larger than that in the SM, and could be detected in the experiment. With a relative small $f$ value ($\sim 500$ GeV) and large $\sqrt{s}$ ($\sim 1$ TeV), the LHT could enhance the SM cross section by three orders and reach 0.1 fb, which might be possible to be seen at future linear colliders. While the LHT have much smaller effect on process $e^+e^- \to b\bar{s}H$ and $pp \to b\bar{s}$ thus couldn’t be detected at future linear colliders and LHT.

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References

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

[2] F. Gabbiani and A. Masiero, Nucl. Phys. B 322, 235 (1989); F. Gabbiani,
E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B 477, 321 (1996) [arXiv:hep-ph/9604387]; M. Misiak, S. Pokorski and J. Rosiek, Adv. Ser. Direct. High Energy Phys. 15, 795 (1998) [arXiv:hep-ph/9703442].

[3] N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Rev. Lett. 86, 4757 (2001) [arXiv:hep-th/0104005]; N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Lett. B 513, 232 (2001) [arXiv:hep-ph/0105239].

[4] For reviews and comprehensive collection of references, see:

M. Schmaltz, and D. Tucker-Smith, Ann. Rev. Nucl. Part. Sci. 55, 229-270 (2005) [arXiv:hep-ph/0502182]; M. Perelstein, Prog. Part. Nucl. Phys. 58, 247-291 (2007) [arXiv:hep-ph/0512128].

[5] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson, JHEP 0207, 034 (2002) [arXiv:hep-ph/0206021].

[6] T. Han, H. E. Logan, B. McElrath, and L. T. Wang, Phys. Rev. D 67, 095004 (2003) [arXiv:hep-ph/0301040].

[7] C. Csaki, J. Hubisz, G.D. Kribs, P. Meade, and J. Terning, Phys. Rev. D 67, 115002 (2003) [arXiv:hep-ph/0211124]; ibid, Phys. Rev. D 68, 035009 (2003) [arXiv:hep-ph/0303236]; J. L. Hewett, F. J. Petriello, and T. G. Rizzo, JHEP 0310, 062 (2003) [arXiv:hep-ph/0211218]; Mu-Chun Chen and Sally Dawson, Phys. Rev. D 70, 015003 (2004) [arXiv:hep-ph/0311032]; W. Kilian and J. Reuter, Phys. Rev. D 70, 015004 (2004) [arXiv:hep-ph/0311095]; Zhenyu Han and Witold Skiba, Phys. Rev. D 72, 035005 (2005) [arXiv:hep-ph/0506206].

[8] H. C. Cheng and I. Low, JHEP0309, 051 (2003) [arXiv:hep-ph/0308199]; JHEP0408, 061 (2004) [arXiv:hep-ph/0405243].

[9] I. Low, JHEP0410, 067 (2004) [arXiv:hep-ph/0409025].

[10] C. S. Huang, X. H. Wu and S. H. Zhu, J. Phys. G 25, 2215 (1999) [arXiv:hep-ph/9902474].
[11] B. Mele, S. Petrarca and A. Soddu, Phys. Lett. B 435 (1998) 401 [arXiv:hep-ph/9805498].

[12] J. M. Yang and C. S. Li, Phys. Rev. D 49, 3412 (1994) [Erratum-ibid. D 51, 3974 (1995)]; J. I. Guasch, arXiv:hep-ph/9906517; J. J. Cao, G. Eilam, M. Frank, K. Hikasa, G. L. Liu, I. Turan and J. M. Yang, Phys. Rev. D 75, 075021 (2007) [arXiv:hep-ph/0702264].

[13] F. Tabbakh, J. J. Liu, W. G. Ma, R. Y. Zhang and H. S. Hou, Commun. Theor. Phys. 44 (2005) 651.

[14] J. Hubisz and P. Meade, Phys. Rev. D 71, 035016 (2005) [arXiv:hep-ph/0411264].

[15] J. Hubisz, S. J. Lee and G. Paz, JHEP 0606, 041 (2006) [arXiv:hep-ph/0512169].

[16] M. Blanke, A. J. Buras, A. Poschenrieder, C. Tarantino, S. Uhlig and A. Weiler, JHEP 0612, 003 (2006) [arXiv:hep-ph/0605214].

[17] M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig and A. Weiler, JHEP 0701, 066 (2007) [arXiv:hep-ph/0610298].

[18] M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig and A. Weiler, Phys. Lett. B 646, 253 (2007) [arXiv:hep-ph/0609284].

[19] M. Blanke, A. J. Buras, B. Duling, S. Recksiegel and C. Tarantino, Acta Phys. Polon. B 41, 657 (2010) [arXiv:0906.5454 [hep-ph]].

[20] J. Kublbeck, M. Bohm and A. Denner, Comput. Phys. Commun. 60, 165 (1990); T. Hahn, Comput. Phys. Commun. 140, 418 (2001) [arXiv:hep-ph/0012260].

[21] G. Passarino and M. J. G. Veltman, Nucl. Phys. B 160, 151 (1979).

[22] G. J. van Oldenborgh and J. A. M. Vermaseren, Z. Phys. C 46, 425 (1990); A. Denner, Fortsch. Phys. 41, 307 (1993) [arXiv:0709.1075 [hep-ph]].
[23] A. J. Buras, A. Poschenrieder, S. Uhlig and W. A. Bardeen, JHEP 0611, 062 (2006) [arXiv:hep-ph/0607189].

[24] K. Nakamura et al. [Particle Data Group], J. Phys. G 37, 075021 (2010).

[25] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) [arXiv:1207.7214 [hep-ex]].

[26] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[27] O. Mena and S. J. Parke, Phys. Rev. D 69, 117301 (2004) [arXiv:hep-ph/0312131].

[28] R. N. Mohapatra et al., Rept. Prog. Phys. 70, 1757 (2007) [arXiv:hep-ph/0510213].

[29] G. Ahuja, M. Gupta and M. Randhawa, arXiv:hep-ph/0611324.

[30] F. P. An et al. [DAYA-BAY Collaboration], Phys. Rev. Lett. 108, 171803 (2012) [arXiv:1203.1669 [hep-ex]].

[31] J. Hubisz, P. Meade, A. Noble and M. Perelstein, JHEP 0601, 135 (2006) [arXiv:hep-ph/0506042].

[32] M. Blanke, A. J. Buras, S. Recksiegel, C. Tarantino and S. Uhlig, Phys. Lett. B 657, 81 (2007) [hep-ph/0703254].

[33] M. Blanke, A. J. Buras, S. Recksiegel and C. Tarantino, arXiv:0805.4393 [hep-ph].

[34] I. I. Bigi, M. Blanke, A. J. Buras and S. Recksiegel, JHEP 0907, 097 (2009) [arXiv:0904.1545 [hep-ph]].

[35] K. Blum, Y. Grossman, Y. Nir and G. Perez, Phys. Rev. Lett. 102, 211802 (2009) [arXiv:0903.2118 [hep-ph]].

[36] X. F. Han, L. Wang and J. M. Yang, Phys. Rev. D 78, 075017 (2008) [arXiv:0807.4480 [hep-ph]].
[37] M. Blanke, A. J. Buras, B. Duling, A. Poschenrieder and C. Tarantino, JHEP 0705, 013 (2007) [arXiv:hep-ph/0702136].