$R_b$ Constraints on Littlest Higgs Model with T-parity

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Abstract

In the framework of the littlest Higgs model with T-parity (LHT), we study the contributions of the T-even and T-odd particles to the branching ratio $R_b$. We find that the precision data of $R_b$ can give strong constraints on the masses of T-odd fermions.

PACS numbers: 14.80.Cp,12.60.Fr,11.30.Qc
I. INTRODUCTION

The little Higgs theory was proposed [1] as a possible solution to the hierarchy problem and so far remains a popular candidate for new physics beyond the SM. The littlest Higgs model [2] is a cute economical implementation of little Higgs, but is found to be subject to strong constraints from electroweak precision tests [3], which would require raising the mass scale of the new particles to far above TeV scale and thus reintroduce the fine-tuning in the Higgs potential [4]. To tackle this problem, a discrete symmetry called T-parity is proposed [5], which forbids the tree-level contributions from the heavy gauge bosons to the observables involving only SM particles as external states. With the running of the LHC, these little Higgs models will soon be put to the test. Since these little Higgs models mainly alter the properties of the Higgs boson and the top quark, hints of these models may be unraveled from various Higgs boson and top quark processes [6].

The branching ratio $R_b$ is defined as

$$R_b \equiv \frac{\Gamma(Z \to b\bar{b})}{\Gamma(Z \to \text{hadrons})},$$

which can provide a precision test of the SM and a sensitive probe of new physics [7]. In the SM most of the electroweak oblique and QCD corrections cancel between numerator and denominator, and the non-decoupling top quark loop effects in the $Zb\bar{b}$ vertex offer a possibility of bounding the top quark mass. In the LHT there are new heavy mirror quarks interacting with gauge bosons, which can contribute to the $R_b$. Therefore, it is possible to give some constraints on the mirror quark masses via their radiative corrections to $R_b$.

The contributions of the LHT to $R_b$ was firstly discussed in [8], which, however, only considered the contributions from the diagrams involving the exchange of the SM Goldstone boson $\pi^\pm$ and neglected the mirror quark contributions under the assumption of flavor-diagonal and flavor-independent mirror quark Yukawa couplings. In this paper, we consider the general situation and examine the contributions of both T-even and T-odd particles to the $R_b$.

The work is organized as follows. In Sec. II we recapitulate the LHT model and discuss the new flavor interactions which will contribute to the decay $Z \to b\bar{b}$. In Sec. III we calculate the one-loop contributions of the LHT to the branching ratio $R_b$ and present constraint of $R_b$ on the mirror quark masses. Finally, we give our conclusions in Sec. IV.
II. THE LITTLEST HIGGS MODEL WITH T-PARITY

The LHT model is based on a non-linear sigma model describing the spontaneous breaking of a global $SU(5)$ down to a global $SO(5)$ by a $5 \times 5$ symmetric tensor at the scale $f \sim \mathcal{O}(\text{TeV})$. From the $SU(5)/SO(5)$ breaking, there arise 14 Goldstone bosons which are described by the "pion" matrix $\Pi$, given explicitly by

$$
\Pi = \begin{pmatrix}
-\omega^0 \sqrt{2} & -\eta \sqrt{20} & -i \omega^+ \sqrt{2} & -i \phi^+ \sqrt{2} & -i \phi^0 \sqrt{2} \\
\omega^- \sqrt{2} & \omega^0 \sqrt{2} & \omega^+ \sqrt{2} & \phi^+ \sqrt{2} & \phi^0 \sqrt{2} \\
-i \pi^- \sqrt{2} & \frac{v+h+\pi^0}{\sqrt{2}} & \sqrt{4/5}\eta & -i \pi^0 \sqrt{2} & \frac{v+h+\pi^0}{\sqrt{2}} \\
i \phi^- \sqrt{2} & \frac{i \pi^- + \phi^0}{\sqrt{2}} & \frac{i \pi^0}{\sqrt{2}} & \frac{\omega^- \sqrt{2}}{\sqrt{2}} & \frac{\omega^0 \sqrt{2}}{\sqrt{2}} \\
i \phi^- \sqrt{2} & \frac{i \omega^0 + \phi^0}{\sqrt{2}} & \frac{v+h+\pi^0}{\sqrt{2}} & \frac{\omega^+ \sqrt{2}}{\sqrt{2}} & \frac{\omega^0 \sqrt{2}}{\sqrt{2}}
\end{pmatrix}.
$$

Under T-parity the SM Higgs doublet $H = (-i\pi^+/\sqrt{2}, (v+h+i\pi^0)/2)^T$ is T-even while other fields are T-odd. A subgroup $[SU(2) \times U(1)]_1 \times [SU(2) \times U(1)]_2$ of the $SU(5)$ is gauged and at the scale $f$ it is broken into the SM electroweak symmetry $SU(2)_L \times U(1)_Y$. The Goldstone bosons $\omega^0, \omega^\pm$ and $\eta$ are respectively eaten by the new T-odd gauge bosons $Z_H, W_H$ and $A_H$, which obtain masses at $\mathcal{O}(v^2/f^2)$

$$M_{W_H} = M_{Z_H} = fg \left(1 - \frac{v^2}{8f^2}\right), \quad M_{A_H} = \frac{fg'}{\sqrt{5}} \left(1 - \frac{5v^2}{8f^2}\right),$$

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

The Goldstone bosons $\pi^0$ and $\pi^\pm$ are eaten by the T-even $Z$ and $W$ bosons of the SM, which obtain masses at $\mathcal{O}(v^2/f^2)$

$$M_{W_L} = \frac{gv}{2} \left(1 - \frac{v^2}{12f^2}\right), \quad M_{Z_L} = \frac{gv}{2 \cos \theta_W} \left(1 - \frac{v^2}{12f^2}\right).$$

The photon $A_L$ is also T-even and remains massless.

For each SM quark, a copy of mirror quark with T-odd quantum number is added in order to preserve the T-parity. We denote them by $u^i_H$ and $d^i_H$, where $i = 1, 2, 3$ are the generation index. In $\mathcal{O}(v^2/f^2)$ their masses are given by

$$m_{d^i_H} = \sqrt{2}\kappa_{q^i} f, \quad m_{u^i_H} = m_{d^i_H} \left(1 - \frac{v^2}{8f^2}\right),$$

where $\kappa_{q^i}$ are the diagonalized Yukawa couplings of the mirror quarks.
Note that new flavor interactions arise between the mirror fermions and the SM fermions, mediated by the T-odd gauge bosons or T-odd Goldstone bosons. In general, besides the charged-current flavor-changing interactions, the FCNC interactions between the mirror fermions and the SM fermions can also arise from the mismatch of rotation matrices. For example, there exist FCNC interactions between the mirror up-type (down-type) quarks and the SM up-type (down-type) quarks, where the mismatched mixing matrix is denoted by $V_{H_u} (V_{H_d})$ with $V_{H_u}^\dagger V_{H_d} = V_{CKM}$. We follow [9] to parameterize $V_{H_d}$ with three angles $\theta_{12}^d, \theta_{23}^d, \theta_{13}^d$ and three phases $\delta_{12}^d, \delta_{23}^d, \delta_{13}^d$

\[
\begin{pmatrix}
  c_{12}^d c_{13}^d & s_{12}^d c_{13}^d e^{-i\delta_{12}^d} & s_{13}^d e^{-i\delta_{13}^d} \\
  -s_{12}^d c_{13}^d e^{i\delta_{12}^d} - c_{12}^d s_{13}^d e^{i(\delta_{13}^d - \delta_{23}^d)} & c_{12}^d s_{23}^d - s_{12}^d s_{23}^d e^{i(\delta_{13}^d - \delta_{23}^d)} & s_{13}^d e^{-i\delta_{13}^d} \\
  s_{12}^d s_{23}^d e^{i(\delta_{13}^d + \delta_{23}^d)} - c_{12}^d s_{13}^d e^{i\delta_{13}^d} & -c_{12}^d s_{23}^d e^{i\delta_{23}^d} - s_{12}^d s_{23}^d e^{i(\delta_{13}^d - \delta_{12}^d)} & c_{13}^d.
\end{pmatrix}
\]  

(6)

III. $R_b$ in the LHT Model

Fig. [1] shows the Feynman diagrams via which LHT gives the corrections to $\Gamma(Z \to b\bar{b})$. The corrections are from both T-even and T-odd particles. The contributions of T-even particles are from the modified coupling $Z t \bar{t}$, $W t \bar{b}$ and $\pi^+ t \bar{b}$, and loops involving the top quark T-even partner (T-quark). The diagrams of T-odd particles are induced by the interactions between the SM quarks and the mirror quarks mediated by the heavy T-odd gauge bosons or Goldstone bosons. The corrections of LHT to the $\Gamma(Z \to d\bar{d})$ and $\Gamma(Z \to s\bar{s})$ are similar to $\Gamma(Z \to b\bar{b})$. For the $\Gamma(Z \to u\bar{u})$ and $\Gamma(Z \to c\bar{c})$, the corrections are only from the T-odd particles, and corrections from the T-even particle can be neglected safely due to the small coupling of $Z T \bar{u}$ and $Z T \bar{c}$. In this work, our purpose is to examine the $R_b$ dependence on the mirror quarks mass, and adopt the method of Bernabeu, Pich, and Santamaria (BPS) to calculate various hadronic decay widths of $Z$ boson [10, 11, 12]. In Appendix, we present the calculation in detail.

In LHT, the branching ratio of $Z \to b\bar{b}$ can be expressed as

\[
R_b \simeq R_b^{SM} (1 + \frac{\delta \Gamma_b}{\Gamma_b^{SM}} - R_b^{SM} \frac{\delta \Gamma_{had}}{\Gamma_b^{SM}}),
\]

(7)

where $R_b^{SM}$ and $\Gamma_b^{SM}$ are the SM predictions for the branching ratio of $Z \to b\bar{b}$ and the width $\Gamma(Z \to b\bar{b})$, $\delta \Gamma_b$ and $\delta \Gamma_{had}$ are the correction of LHT to the $\Gamma_b^{SM}$ and $\Gamma^{SM}(Z \to hadrons)$, respectively.
In the numerical calculations we take the Fermi constant \( G_F \), the fine-structure constant \( \alpha_{M_Z} \), \( Z \)-boson mass \( M_{Z_L} \), fermion masses \( m_f \), and the electroweak mixing angle \( s_W = \sin \theta_W \) as input parameters. The LHT parameters relevant to our calculation are the scale \( f \), the ratio between top quark Yukawa couplings \( r = \frac{\lambda_t}{\lambda_2} \), the mirror quark masses and parameters in the matrices \( V_{H_u} \) and \( V_{H_d} \). \( f \) may be as low as 500 GeV, and \( r \) is taken typical value as 1. For the mirror quark masses, from Eq. (5) we get \( m_{u_H} = m_{d_H} \) at \( \mathcal{O}(v/f) \) and further we assume

\[
m_{u_H} = m_{u_H}^2 = m_{d_H} = m_{d_H}^2 \equiv M_{12}, \quad m_{u_H}^3 = m_{d_H}^3 \equiv M_3. \tag{8}
\]

For the matrices \( V_{H_u} \) and \( V_{H_d} \), considering the constraints in [14], we follow them to consider the following four scenarios:

(I) \( V_{H_u} = 1, V_{H_d} = V_{CKM} \).

(II) \( V_{H_d} = 1, V_{H_u} = V_{CKM}^T \).

(III) \( s_{13}^d = 0.5, s_{12}^d = s_{23}^d = 0, \delta_{13}^d = \delta_{13}^{SM}, s_{ij} = s_{ij}^{SM} \) otherwise.

(IV) \( s_{13}^d = 0.5, s_{12}^d = 0.7, s_{23}^d = 0.4, \delta_{12}^d = \delta_{23}^d = 0, \delta_{13}^d = \delta_{13}^{SM} \).

In Figs. 2-5, we plot the branching ratio \( R_b \) versus the first two mirror quark mass \( M_{12} \) for the scenario I, II, III and IV, respectively. The Figs. 2-5 show \( R_b \) can give strict lower bound and upper bound of the first two mirror quark mass for the \( f \) and \( M_3 \) taken. The
FIG. 2: The branching ratio $R_b$ versus the mass of first two family mirror quarks in scenario-I with $f = 500$ GeV, 1000 GeV and 2000 GeV, respectively.

FIG. 3: Same as Fig. 2, but for scenario-II.

constraints are sensitive to the scale $f$, and the allowed regions of $M_{12}$ become larger with the increasing $f$. Further, the $R_b$ favors a large value of $M_{12}$ for a large value of $f$.

In scenario I and scenario II, the up-type Yukawa interactions and the down-type quark Yukawa interactions are diagonal, respectively. However, scenario III and scenario IV are two large mixing scenarios, and the angle $s_{13}^d$ is set large so that the third generation mass dependence can be more sensitive. For example, when $f = 1$ TeV (2 TeV), the four lines in Fig. 2 and Fig. 3 are almost overlapped for scenario I and scenario II, and this situation can be relaxed for scenario III and scenario IV. Besides, for $f = 500$ GeV, $M_3 = 3000$ GeV can be allowed in scenario I and scenario II, but be ruled out in scenario III and scenario IV by the $2\sigma$ $R_b$ constraints.
IV. CONCLUSION

In the framework of littlest Higgs model with T-parity, we studied the loop contributions of the T-even and T-odd particles to the branching ratio $R_b$ for four different scenarios. We found that the precision measurement data of $R_b$ can give strong constraints on the mirror quark masses. For the values of $f$ and $M_3$ in various scenarios of $V_{Hd}$, $R_b$ can give strict lower bound and upper bound for the mass $M_{12}$ of the first two generations of mirror quarks, and the allowed regions of $M_{12}$ become larger as $f$ gets large. Further, the $R_b$ data favors a large value of $M_{12}$ in case of a large $f$. Besides, the $R_b$ constraints on the masses of three generation mirror quarks depend on the texture of $V_{Hd}$, and are more sensitive to the mass $M_3$ of the third generation of mirror quarks in scenarios III and IV than in scenarios I and II. For example, when $f = 500$ GeV, $M_3 = 3000$ GeV is allowed in scenarios I and II, but ruled out in scenarios III and IV by the 2σ $R_b$ constraints.
Acknowledgement

We thank L. Wang, J. M. Yang and C. P. Yuan for discussions. This work was supported by the National Natural Science Foundation of China (NNSFC) under Nos. 10821504, 10725526 and 10635030.

APPENDIX A: THE HEAVY QUARK LOOP CONTRIBUTIONS TO $Z \to b\bar{b}$

According to the BPS method [10, 11, 12], we give the expressions of hadronic decay widths of $Z$-boson:

$$\Gamma_q = \frac{3m_Z}{12\pi} (v_q^2 + a_q^2) \left[ 1 + \frac{\alpha}{\pi} \left( \frac{Z^q_L}{(Z^q_L)^2 + (Z^q_R)^2} \right) F_q \right] \quad (q = b, d, s, u, c),$$  \hfill (A1)

where $v_q = \frac{Z^q_L + Z^q_R}{2}$ and $a_q = \frac{Z^q_L - Z^q_R}{2}$ with $Z^q_L$ and $Z^q_R$ being the left- and right-handed couplings of $Zq\bar{q}$, respectively.

$$F_{b,d,s} = V_{cha}(t, W, \pi) + V_{cha}(T, W, \pi) + V_{cha}(u_H^i, W_H, \omega) + V_{neu}(d_H^i, Z_H, \omega^0) + V_{neu}(d_H^i, A_H, \eta) + V_{mix}(t, T, W, \pi),$$

$$F_{u,c} = V_{cha}(d_H^i, W_H, \omega) + V_{neu}(u_H^i, Z_H, \omega^0) + V_{neu}(u_H^i, A_H, \eta),$$ \hfill (A2)

where

$$V_{cha}(f, V, S) = F^{(a)} + F^{(b)} + F^{(c)+(d)} + F^{(e)+(f)} + F^{(g)} + F^{(i)+(j)},$$

$$V_{neu}(f, V, S) = F^{(a)} + F^{(c)+(d)} + F^{(g)} + F^{(i)+(j)},$$

$$V_{mix}(f, f', V, S) = F^{(k)+(l)} + F^{(m)+(n)}. \hfill (A3)$$
The above equations are the corresponding explicit expressions of the Feynman diagrams in Fig. 11, which are given by

\[
F^{(a)} = -\frac{1}{g^2 s_w^2} \left| c_3 \right|^2 \left\{ \frac{r}{2} \left[ \frac{r(r - 2)}{(r - 1)^2} \ln r + \frac{r}{r - 1} \right] + \frac{r}{2} \left[ \frac{r(r - 2)}{(r - 1)^2} \ln r - \frac{r}{r - 1} \right] \right\},
\]

\[
F^{(b)} = -\frac{3}{2g^2 s_w^2} \left| c_3 \right|^2 g_{ZV} \left( \frac{r^2}{(r - 1)^2} \ln r - \frac{r}{r - 1} \right),
\]

\[
F^{(c)+(d)} = \frac{1}{g^2 s_w^2} \left| c_3 \right|^2 g_{ZL} \left[ \frac{r^2}{(r - 1)^2} \ln r - \frac{r}{r - 1} \right],
\]

\[
F^{(e)+(f)} = \frac{1}{g^2 s_w^2} \left| c_3 \right|^2 g_{ZS} \left[ \frac{r}{(r - 1)^2} \ln r - \frac{1}{r - 1} \right],
\]

\[
F^{(g)} = -\frac{1}{2g^2 s_w^2} \left| a_3 \right|^2 \left\{ \frac{r^2}{2} \left[ \frac{r(r - 2)}{(r - 1)^2} \ln r + \frac{2r - 1}{r - 1} \right] + \frac{r^2}{2} \left[ \frac{r(r - 2)}{(r - 1)^2} \ln r - \frac{r}{r - 1} \right] \right\},
\]

\[
F^{(h)} = \frac{1}{4g^2 s_w^2} \left| a_3 \right|^2 g_{ZSV} \left[ \frac{r^2}{(r - 1)^2} \ln r - \frac{r}{r - 1} \right],
\]

\[
F^{(i)+(j)} = -\frac{2}{g^2 s_w^2} \left| t^c \right|^2 g_{ZL} \left[ \frac{r^2}{2} \left[ \frac{r}{r - 1} \ln r - \frac{r}{r - 1} \right] \right] - \frac{Z_L^{TT}}{\sqrt{r r'}} \frac{1}{r - 1} \left[ \frac{r'}{r' - 1} \ln r' - \frac{r}{r - 1} \ln r \right],
\]

\[
F^{(m)+(n)} = \frac{1}{2g^2 s_w^2} \left| a_3 \right|^2 \left\{ \frac{2Z_L^{TT}}{r'} \left[ \frac{r'}{r' - 1} \ln r' - \frac{r}{r - 1} \ln r \right] - \frac{Z_R^{TT}}{\sqrt{r r'}} \right\} \left( \Delta + 1 + \frac{1}{r - 1} \left[ \frac{r^2}{r - 1} \ln r' - \frac{r^2}{r - 1} \ln r \right] \right),
\]

with

\[
\Delta \equiv \frac{2}{n - 4} + \gamma + \ln (m_V^2 / 4\pi \mu^2) - \frac{3}{2},
\]

\[
r = m_f^2 / m_V^2, \quad r' = m_{f'}^2 / m_V^2.
\]

The coupling constant appearing above are from

\[
V \tilde{f} q : i \gamma^\mu (c_3^T P_L + d_3^T P_R), \quad Z \tilde{f} f : i \gamma^\mu (Z_L^T P_L + Z_R^T P_R), \quad Z \tilde{f} \bar{f} : i \gamma^\mu (Z_L^T P_L + Z_R^T P_R),
\]

\[
Z S^+ S^- : i g_{VSS} (p_{S+}^\mu - p_{S-}^\mu), \quad \bar{Z} V^+ V^- : g_{ZV} g_{SS}^{\mu \nu}, \quad Z^\rho V^{\mu} V^{-\nu} : i g_{ZV V} \left[ g^{\mu \nu} (p_+ - p_-) + g^{\nu \rho} (p_- - p_Z) + g^{\rho \mu} (p_Z - p_+) \right],
\]

where \( f, V \) and \( S \) represent fermion, gauge bosons and scalar particles involved in the loops, respectively. The explicit expressions of these parameters are complicated at \( O(v^2 / f^2) \) and
can be found in [15, 16].

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