Top quark decay to a 125 GeV Higgs in BLMSSM

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Abstract

In this paper, we calculate the top quark rare decay $t \rightarrow ch$ in a supersymmetric extension of the standard model where baryon and lepton numbers are local gauge symmetries. Adopting reasonable assumptions on the parameter space, we find that the branching ratios of $t \rightarrow ch$ can reach $10^{-3}$, which can be detected in near future.

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I. INTRODUCTION

Top quark plays a special role in the standard model (SM) and holds great promise in revealing the secret of new physics beyond the SM. The running LHC is a top-quark factory, and provides a great opportunity to seek out top-quark rare decays. Among those rare processes, the flavor-changing neutral current (FCNC) decays \( t \rightarrow ch \) deserve special attention, since the branching ratios (BRs) of those rare processes are strongly suppressed in the SM. In addition, ATLAS and CMS have reported significant excess events which are interpreted probably to be related to the neutral Higgs with mass \( m_{h_0} \sim 124−126 \text{ GeV} \) \[1, 2\]. This implies that the Higgs mechanism to break electroweak symmetry possibly has a solid experimental cornerstone.

In the framework of the SM, the possibility of detecting FCNC decays \( t \rightarrow ch \) is essentially hopeless, since tree level FCNC involving the quarks are forbidden by the gauge symmetries and particle content \[3, 4\]. In particular, it has recently been recognized that the BRs of the process is much smaller \[5, 6\] than originally thought \[7\] which is less then \( 10^{-13} \). In extensions of the SM, the BRs for FCNC top decays can be orders of magnitude larger. For example, the authors of Ref. \[8\], study \( t \rightarrow ch \) process in the framework of the minimal supersymmetric extension of the standard model (MSSM) include the leading set of supersymmetric QCD and supersymmetric electroweak contributions, and get \( \text{Br}^{\text{SUSY-EW}}(t \rightarrow ch) \sim 10^{-8} \), \( \text{Br}^{\text{SUSY-QCD}}(t \rightarrow ch) \sim 10^{-5} \). And a new work about this process in MSSM is discussed in Ref. \[9\], with \( \tan \beta = 1.5 \) or 35 and the mass of SUSY particles about 1 or 2TeV scale, the authors get the branching ratio of \( t \rightarrow ch \) can only reach \( 3 \times 10^{-6} \), which is much smaller than previous results obtained before the advent of the LHC.

Physicists have been interested in the MSSM \[10-13\] for a long time. However, since the matter-antimatter is asymmetry in the universe, baryon number (B) should be broken. On the other hand, since heavy majorana neutrinos contained in the seesaw mechanism can induce the tiny neutrino masses \[14, 15\] to explain the neutrino oscillation experiment, so the lepton number (L) is also expected to be broken. A minimal supersymmetric extension...
of the SM with local gauged B and L (BLMSSM) is more favorite \[16, 17\]. Since the new quarks are vector-like with respect to the strong, weak and electromagnetic interactions to cancel anomalies, one obtains that their masses can be above 500 GeV without assuming large couplings to the Higgs doublets in this model. Therefore, there are no Landau poles for the Yukawa couplings here.

In BLMSSM, B and L are spontaneously broken near the weak scale, the proton decay is forbidden, and the three neutrinos get mass from the extended seesaw mechanism at tree level \[3, 4, 16, 17\]. Therefore, the desert between the grand unified scale and the electroweak scale is not necessary, which is the main motivation for the BLMSSM.

The CMS \[18\] and ATLAS \[19\] experiments of LHC have studied many possible signals of the MSSM, and set very strong bounds on the gluino and squarks masses with R-parity conservation. However, in BLMSSM, the predictions and bounds for the collider experiments should be changed \[16, 17, 20\]. In addition, the lepton number violation could be detected at the LHC from the decays of right handed neutrinos \[3, 4, 21\], and we could also look for the baryon number violation in the decays of squarks and gauginos \[22\]. Since there are some exotic fields, and exist couplings between exotic quark fields and SM quarks in the superpotential, so it will cause flavor changing processes, and the BRs for FCNC top decays can be orders of magnitude larger.

In this paper we analyze the corrections to the top-quark decay $t \to ch$ in BLMSSM. This paper is composed of the sections as follows. In section II, we present the main ingredients of the BLMSSM. In section III, we present the theoretical calculation on the $t \to ch$ processes. Section IV is devoted to the numerical analysis. Our conclusions are summarized in Section V.

II. A SUPERSYMMETRIC EXTENSION OF THE SM WHERE B AND L ARE LOCAL GAUGE SYMMETRIES

The local gauge B and L is base on the gauge group: $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L$. In BLMSSM, to cancel the $B$ and $L$ anomalies, the exotic superfields should
include the new quarks $\hat{Q}_4$, $\hat{U}_c^c$, $\hat{D}_c^c$, $\hat{Q}_5$, $\hat{U}_5$, $\hat{D}_5$, and the new leptons $\hat{L}_4$, $\hat{E}_c^c$, $\hat{N}_4^c$, $\hat{L}_5^c$, $\hat{E}_5$, $\hat{N}_5$. In addition, the new Higgs chiral superfields $\hat{\Phi}_B$ and $\hat{\phi}_B$ acquire nonzero vacuum expectation values (VEVs) to break baryon number spontaneously, the superfields $\hat{\Phi}_L$ and $\hat{\phi}_L$ acquire nonzero VEVs to break lepton number spontaneously. The model also introduce the superfields $\hat{X}$, $\hat{X}'$ to avoid stability for the exotic quarks. Actually, the lightest superfields can be a candidate for dark matter. The properties of these superfields in BLMSSM are summarized in Table I.

**TABLE I: The properties of superfields in BLMSSM**

| superfield | $SU(3)$ | $SU(2)$ | $U(1)_Y$ | $U(1)_B$ | $U(1)_L$ |
|------------|---------|---------|----------|----------|----------|
| $\hat{Q}_4$ | 3       | 2       | 1/6      | $B_4$    | 0        |
| $\hat{U}_c^c$ | 3       | 1       | -2/3     | $-B_4$   | 0        |
| $\hat{D}_c^c$ | 3       | 1       | 1/3      | $-B_4$   | 0        |
| $\hat{Q}_5$ | 3       | 2       | -1/6     | $-(1 + B_4)$ | 0     |
| $\hat{U}_5$ | 3       | 1       | 2/3      | $1 + B_4$ | 0        |
| $\hat{D}_5$ | 3       | 1       | -1/3     | $1 + B_4$ | 0        |
| $\hat{L}_4$ | 1       | 2       | -1/2     | 0        | $L_4$    |
| $\hat{E}_c^c$ | 1       | 1       | 1        | 0        | $-L_4$   |
| $\hat{N}_4^c$ | 1       | 1       | 0        | 0        | $-L_4$   |
| $\hat{L}_5^c$ | 1       | 2       | 1/2      | 0        | $-(3 + L_4)$ |
| $\hat{E}_5$ | 1       | 1       | -1       | 0        | $3 + L_4$ |
| $\hat{N}_5$ | 1       | 1       | 0        | 0        | $3 + L_4$ |
| $\hat{\Phi}_B$ | 1       | 1       | 0        | 1        | 0        |
| $\hat{\phi}_B$ | 1       | 1       | 0        | -1       | 0        |
| $\hat{\Phi}_L$ | 1       | 1       | 0        | 0        | $-2$     |
| $\hat{\phi}_L$ | 1       | 1       | 0        | 0        | 2        |
| $\hat{X}$ | 1       | 1       | 0        | $2/3 + B_4$ | 0    |
| $\hat{X}'$ | 1       | 1       | 0        | $-(2/3 + B_4)$ | 0   |
In BLMSSM, the super potential is written as

$$\mathcal{W}_{BLMSSM} = \mathcal{W}_{MSSM} + \mathcal{W}_B + \mathcal{W}_L + \mathcal{W}_X,$$  

(1)

where $\mathcal{W}_{MSSM}$ is superpotential of the MSSM, and the concrete form of $\mathcal{W}_B$, $\mathcal{W}_L$ and $\mathcal{W}_X$ are

$$\mathcal{W}_B = \lambda_Q \tilde{Q}_4 \tilde{Q}_4^c \tilde{\Phi}_B + \lambda_U \tilde{U}_4 \tilde{U}_5 \tilde{\Phi}_B + \lambda_D \tilde{D}_4 \tilde{D}_5 \tilde{\Phi}_B + \mu_B \tilde{\Phi}_B \tilde{\Phi}_B$$

$$+ Y_{u_4} \tilde{Q}_4 \tilde{H}_u \tilde{U}_4^c + Y_{d_4} \tilde{Q}_4 \tilde{H}_d \tilde{D}_4^c + Y_{u_5} \tilde{Q}_5 \tilde{H}_u \tilde{U}_5 + Y_{d_5} \tilde{Q}_5 \tilde{H}_d \tilde{D}_5^c,$$

$$\mathcal{W}_L = Y_{e_4} \tilde{L}_4 \tilde{H}_u \tilde{E}_4^c + Y_{\nu_4} \tilde{L}_4 \tilde{H}_u \tilde{\nu}_4^c + Y_{\nu_5} \tilde{L}_5 \tilde{H}_u \tilde{\nu}_5 + Y_{\nu_5} \tilde{L}_5 \tilde{H}_d \tilde{\nu}_5$$

$$+ Y_{\nu_5} \tilde{L}_5 \tilde{H}_d \tilde{\nu}_5 + \lambda_{\nu_5} \tilde{\nu}_5 \tilde{\nu}_5 \tilde{\Phi}_L + \mu_L \tilde{\Phi}_L \tilde{\Phi}_L,$$

$$\mathcal{W}_X = \lambda_1 \tilde{Q}_5 \tilde{Q}_5^c \tilde{X} + \lambda_2 \tilde{U}_5 \tilde{U}_5^c \tilde{X}' + \lambda_3 \tilde{D}_5 \tilde{D}_5^c \tilde{X}' + \mu_X \tilde{X} \tilde{X}'$$,  

(2)

and we could see that since $\mathcal{W}_X$ contains superfields $X$ and $Q_5$ ($U_5$, $D_5$ and $X'$) couple to all generations of SM quarks, so FCNC processes can be generated.

Correspondingly, the soft breaking terms $\mathcal{L}_{soft}$ are generally given as

$$\mathcal{L}_{soft} = \mathcal{L}_{soft}^{MSSM} - (m_{\nu^c}_{11})^2 \tilde{\nu}_1^c \bar{\tilde{\nu}}_1 - m_{\tilde{Q}_4}^2 \tilde{Q}_4^c \tilde{Q}_4 - m_{\tilde{U}_4}^2 \tilde{U}_4^c \tilde{U}_4 - m_{\tilde{D}_4}^2 \tilde{D}_4^c \tilde{D}_4$$

$$- m_{\tilde{Q}_5}^2 \tilde{Q}_5^c \tilde{Q}_5 - m_{\tilde{U}_5}^2 \tilde{U}_5^c \tilde{U}_5 - m_{\tilde{D}_5}^2 \tilde{D}_5^c \tilde{D}_5 - m_{\tilde{L}_4}^2 \tilde{L}_4^c \tilde{L}_4$$

$$- m_{\tilde{E}_4}^2 \tilde{E}_4^c \tilde{E}_4 - m_{\tilde{\nu}_5}^2 \tilde{\nu}_5^c \tilde{\nu}_5 - m_{\tilde{\nu}_5}^2 \tilde{\nu}_5^c \tilde{\nu}_5 - m_{\tilde{\Phi}_B}^2 \tilde{\Phi}_B \tilde{\Phi}_B$$

$$- m_{\tilde{\nu}_L}^2 \tilde{\nu}_L \tilde{\nu}_L - m_{\tilde{\nu}_L}^2 \tilde{\nu}_L \tilde{\nu}_L - \left(m_B \lambda_B \lambda_B + m_L \lambda_L \lambda_L + h.c.\right)$$

$$+ \left\{ A_{u_4} Y_{u_4} \tilde{Q}_4 \tilde{H}_u \tilde{U}_4^c + A_{d_4} Y_{d_4} \tilde{Q}_4 \tilde{H}_d \tilde{D}_4^c + A_{u_5} Y_{u_5} \tilde{Q}_5 \tilde{H}_u \tilde{U}_5 + A_{d_5} Y_{d_5} \tilde{Q}_5 \tilde{H}_d \tilde{D}_5^c$$

$$+ A_{BQ} \lambda_Q \tilde{Q}_4 \tilde{Q}_5^c \tilde{\Phi}_B + A_{BU} \lambda_U \tilde{U}_4^c \tilde{U}_5 \tilde{\Phi}_B + A_{BD} \lambda_D \tilde{D}_4 \tilde{D}_5 \tilde{\Phi}_B + B_B \mu_B \tilde{\Phi}_B \tilde{\Phi}_B + h.c.\right\}$$

$$+ \left\{ A_{e_5} Y_{e_5} \tilde{L}_5 \tilde{H}_d \tilde{E}_5^c + A_{\nu_5} Y_{\nu_5} \tilde{L}_5 \tilde{H}_d \tilde{\nu}_5^c + A_{e_5} Y_{e_5} \tilde{L}_5 \tilde{H}_d \tilde{E}_5^c + A_{\nu_5} Y_{\nu_5} \tilde{L}_5 \tilde{H}_d \tilde{\nu}_5$$

$$+ A_{\nu_u} \tilde{L}_5 \tilde{H}_u \tilde{\nu}_5 + A_{\nu_u} \lambda_{\nu_u} \tilde{\nu}_5 \tilde{\nu}_5 \tilde{\Phi}_L + B_L \mu_L \tilde{\Phi}_L \tilde{\Phi}_L + h.c.\right\}$$

$$+ \left\{ A_{1} \lambda_1 \tilde{Q}_5 \tilde{Q}_5^c X + A_{2} \lambda_2 \tilde{U}_5 \tilde{U}_5^c X' + A_{3} \lambda_3 \tilde{D}_5 \tilde{D}_5^c X' + B_X \mu_X \tilde{X} \tilde{X}' + h.c.\right\} ,$$  

(3)

with $\mathcal{L}_{soft}^{MSSM}$ representing the soft breaking terms of the MSSM, and $\lambda_B$, $\lambda_L$ are gauginos of $U(1)_B$ and $U(1)_L$, respectively.
To break the local gauge symmetry $SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L$ down to the electromagnetic symmetry $U(1)_e$, the $SU(2)_L$ doublets $H_u$, $H_d$ and the $SU(2)_L$ singlets $\Phi_B$, $\varphi_B$, $\Phi_L$, $\varphi_L$ should obtain the nonzero VEVs $v_u$, $v_d$, $v_B$, and $v_L$, $\overline{v}_L$ respectively.

$$H_u = \begin{pmatrix} H_u^+ \\ \frac{1}{\sqrt{2}}(v_u + H_u^0 + iP^0_u) \end{pmatrix},$$

$$H_d = \begin{pmatrix} H_d^- \\ \frac{1}{\sqrt{2}}(v_d + H_d^0 + iP^0_d) \end{pmatrix},$$

$$\Phi_B = \frac{1}{\sqrt{2}}(v_B + \Phi_B^0 + iP^0_B),$$

$$\varphi_B = \frac{1}{\sqrt{2}}(\overline{v}_B + \varphi_B^0 + i\overline{P}^0_B),$$

$$\Phi_L = \frac{1}{\sqrt{2}}(v_L + \Phi_L^0 + iP^0_L),$$

$$\varphi_L = \frac{1}{\sqrt{2}}(\overline{v}_L + \varphi_L^0 + i\overline{P}^0_L).$$

The mass matrixes of Higgs, exotic quarks and exotic scalar quarks are obtained in our previous work [23], and we list some useful results.

In four-component Dirac spinors, the mass matrix for exotic charged 2/3 quarks is

$$-L_{\mu u}^{mass} = \begin{pmatrix} \tilde{t}_{4R}'' & \tilde{t}_{5R}'' \end{pmatrix} \begin{pmatrix} \frac{1}{\sqrt{2}}Y_{u4}v_u, & \frac{1}{\sqrt{2}} \lambda_{Q4u} \\ -\frac{1}{\sqrt{2}} \lambda_{u4} \overline{v}_B, & \frac{1}{\sqrt{2}} Y_{u5}v_d \end{pmatrix} \begin{pmatrix} t_{4L}'' \\ t_{5L}'' \end{pmatrix} + h.c. \tag{5}$$

and it could be diagonalized by the the unitary transformations

$$\begin{pmatrix} t_{4L}'' \\ t_{5L}'' \end{pmatrix} = U_{\nu}^+ \cdot \begin{pmatrix} t_{4L}' \\ t_{5L}' \end{pmatrix}, \quad \begin{pmatrix} t_{4R}' \\ t_{5R}' \end{pmatrix} = W_{\nu}^+ \cdot \begin{pmatrix} t_{4R}'' \\ t_{5R}'' \end{pmatrix} \tag{6}$$

then we get

$$W_{\nu}^+ \cdot \begin{pmatrix} \frac{1}{\sqrt{2}}Y_{u4}v_u, & \frac{1}{\sqrt{2}} \lambda_{Q4u} \\ -\frac{1}{\sqrt{2}} \lambda_{u4} \overline{v}_B, & \frac{1}{\sqrt{2}} Y_{u5}v_d \end{pmatrix} \cdot U_{\nu} = diag(m_{t4}, m_{t5}) \tag{7}$$

Similarly, The concrete expressions for $4 \times 4$ mass squared matric $M_{\nu}^2$ of exotic charged 2/3 scalar quarks $\tilde{t}^{\nu T} = (\tilde{Q}_4^1, \tilde{U}_4^{cs}, \tilde{Q}_5^{2cs}, \tilde{U}_5)$ are given in appendix B of Ref [23], and it could
be diagonalized by the unitary transformation

$$\tilde{t}''_i = Z^i_j \tilde{t}'_j,$$  \hspace{1cm} (8)

Using the scalar potential and the soft breaking terms, the mass squared matrix for $X, X'$ could be written as

$$-L_{\text{mass}}^{X} = \begin{pmatrix} X^* & X' \end{pmatrix} \begin{pmatrix} \mu^2_X + S_X & -B_X \mu_X \\ -B_X \mu_X & \mu^2_X - S_X \end{pmatrix} \begin{pmatrix} X \\ X' \end{pmatrix},$$  \hspace{1cm} (9)

with $S_X = \frac{g^2}{2}(\frac{v^2}{3} + B_4)(v_B^2 - v_B^2)$. And it could be diagonalized by the unitary transformation $Z_X$

$$Z^\dagger_X \begin{pmatrix} \mu^2_X + S_X & -B_X \mu_X \\ -B_X \mu_X & \mu^2_X - S_X \end{pmatrix} Z_X = \text{diag}(m^2_{X_1}, m^2_{X_2}).$$  \hspace{1cm} (10)

In addition, the four-component Dirac spinor $\tilde{t}X$ is defined as $\tilde{t}X = (\psi_{X}, \bar{\psi}_{X'})^T$, with the mass term $\mu_X \tilde{t}X \tilde{t}X$.

The flavor conservative couplings between the lightest neutral Higgs and charged $2/3$ exotic quarks are

$$\mathcal{L}_{Ht'} = \frac{1}{\sqrt{2}} \sum_{i,j=1}^{2} \left\{ \left[ Y_{u4}(W_t^\dagger)_{i2}(U_t)_{1j} \cos \alpha + Y_{u5}(W_t^\dagger)_{i1}(U_t)_{2j} \sin \alpha \right] h^0 \bar{t}'_i P_L t'_j + \left[ Y_{u4}(U_t^\dagger)_{i1}(W_t)_{2j} \cos \alpha + Y_{u5}(U_t^\dagger)_{i2}(W_t)_{1j} \sin \alpha \right] h^0 \bar{t}'_i P_R t'_j \right\}$$  \hspace{1cm} (11)

with $\alpha$ is defined as

$$\begin{pmatrix} H^0 \\ h^0 \end{pmatrix} = \begin{pmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} H^0_d \\ H^0_u \end{pmatrix},$$  \hspace{1cm} (12)

And the couplings between the lightest neutral Higgs and exotic scalar quarks are

$$\mathcal{L}_{H\tilde{t}'} = \sum_{i,j}^{4} \left\{ \xi^S_{uij} \cos \alpha - \xi^S_{dij} \sin \alpha \right\} h^0 \bar{\tilde{t}}'_i \tilde{t}'_j,$$  \hspace{1cm} (13)

with $\xi^S_{uij}$ and $\xi^S_{dij}$ are defined in appendix C of Ref[23].
In mass basis, we obtain the couplings of quark-exotic quark and the $X$:

$$-\lambda_1(Wt)_i2(ZX)_1jXj\bar{t}_iP_Lu - \lambda_2(U^\dagger t)_i2(ZX)_2jXj\bar{u}P_Lt'_i + h.c. \quad (14)$$

and the couplings between up type quark and the superpartners $\tilde{t}'$, $\tilde{X}$ are

$$-\lambda_1(Z^\dagger t')_i3\bar{t}_i\bar{u}P_L\tilde{X} - \lambda_2(Z^\dagger t')_i4\bar{t}_i\bar{X}P_Lu + h.c. \quad (15)$$

III. THE THEORETICAL CALCULATION ON THE $t \rightarrow ch$ PROCESS

In this section, we present one-loop radiative corrections to the rare decay $t \rightarrow ch$ in BLMSSM. For this process, it is convenient to define an effective interaction vertex:

$$-iT = -ig\bar{c}(p) \left( F_LP_L + F_RP_R \right) t(p') \quad (16)$$

where $p'$ is the momentum of the initial top quark, $p$ is the momentum of the final state charm quark, and form factors $F_L$, $F_R$ are follow from explicit calculation of vertices and mixed self-energies.

The relevant one-loop vertex diagrams of BLMSSM are drawn in Fig.1.

We could see that the FCNC transitions of new physics are mediated by the exotic up type quark $t'$, the neutral scalar particle $X_i$ and there superpartners $\tilde{t}', \tilde{X}$. And the contribution to the form factors could be obtained by direct calculation.
In the equations below, \(m_{\nu'}\), \(m_X\), \(m_{\tilde{p}}\), and \(m_{\tilde{X}}\) denote the mass of the exotic quarks \(t'\), the mass of scalar particle \(X_i\), and the mass of there superpartners \(\tilde{t}'\), \(\tilde{X}\) respectively. \(B_i, C_{ij}\) are the coefficients of the Lorentz-covariant tensors in the standard scalar Passarino-Veltman integrals (Eq.(4.7) in Ref. [25]), and it could be calculated by using 'LoopTools'.

In Fig.1(a), when one-loop diagrams are composed by the neutral scalar particles \(X_i\), and charged \(2/3\) new quarks \(t'\), the contribution to the form factors \(F_L^a\) and \(F_R^a\) are formulated as

\[
F_L^a = \frac{i}{16\pi^2} \sum_{i,j,l} (-a_1 m_c (b_1 h_2 m_{\nu'} C_2 + b_2 h_1 m_{\nu'} (C_0 + C_1 + 2C_2) + 3b_2 h_2 m_{\nu'} C_2)) \\
+ a_2 b_2 (h_1 B_0 + (h_1 m_{\nu'}, h_2 m_{\nu'}, m_{\nu'}) C_0) \\
+ a_2 b_1 m_i (h_2 m_{\nu'} (C_0 + C_1 + C_2) + h_1 m_{\nu'} (C_0 + C_2)) + a_2 b_2 h_1 m_c^2 C_2
\]

\[
F_R^a = \frac{i}{16\pi^2} \sum_{i,j,l} (-a_2 m_c (b_1 h_2 m_{\nu'} (C_0 + C_1 + 2C_2) + b_1 h_1 m_{\nu'} (C_0 + C_2) + b_2 h_1 m_{\nu'} C_2)) \\
+ a_1 b_1 (h_2 B_0 + (h_1 m_{\nu'}, m_{\nu'} + h_2 m_{\nu'}) C_0) \\
+ a_1 b_2 (m_i (h_1 m_{\nu'} (C_0 + C_1 + C_2) + h_2 m_{\nu'} (C_1 + C_2)) + a_1 b_2 h_1 m_c^2 C_2)
\]

with the Passarino-Veltman integrals

\[
B_0 = B_0 (p^2, m_{\nu'}^2, m_{X_i}^2) \\
C_0 = C_0 (p^2, (2p - p')^2, (p - p')^2, m_{\nu'}^2, m_{X_i}^2, m_{\nu'}^2) \\
C_{1,2} = C_{1,2} ((p - p')^2, (2p - p')^2, p^2, m_{\nu'}^2, m_{\nu'}^2, m_{X_i}^2)
\]

and the relevant coefficients are

\[
a_1 = \lambda_1 (U_{\nu'})_{2i} (Z_{X}^i)_{1l}, \quad a_2 = \lambda_2 (U_{\nu'})_{2i} (Z_{X}^i)_{2l}, \\
b_1 = \lambda_2 (U_{\nu'})_{ji} (Z_{X}^j)_{1l}, \quad b_2 = \lambda_1 (U_{\nu'})_{ji} (Z_{X}^j)_{2l}, \\
h_1 = Y_{a_1} (U_{\nu'})_{i1} (W_{\nu'})_{2j} \cos \alpha + Y_{a_2} (U_{\nu'})_{i2} (W_{\nu'})_{1j} \sin \alpha, \\
h_2 = Y_{a_1} (W_{\nu'})_{i2} (U_{\nu'})_{1j} \cos \alpha + Y_{a_2} (W_{\nu'})_{i1} (U_{\nu'})_{2j} \sin \alpha,
\]

In Fig.1(b), when one-loop diagrams are composed by the superpartners \(\tilde{t}'\) and \(\tilde{X}\), \(F_L^b\) and \(F_R^b\) are formulated as
In Fig. 2 we present the relevant self-energy diagrams of the rare decay $t \to ch$ in BLMSSM.

![Diagram of self-energy diagrams](image)

$$F_L^b = \frac{i}{16\pi^2} \sum_{i,j} (a_4 b_4 m_{\tilde{X}} C_0 - a_3 b_4 m_c C_1 - a_4 b_3 m_t C_2)(\cos \alpha \xi_u - \sin \alpha \xi_d)$$

$$F_R^b = \frac{i}{16\pi^2} \sum_{i,j} (a_4 b_4 m_{\tilde{X}} C_0 - a_3 b_4 m_c C_1 - a_4 b_3 m_t C_2)(\cos \alpha \xi_u - \sin \alpha \xi_d)$$

(20)

with

$$C_0 = \lambda_2^*(Z^i_{\tilde{t}'}^i, P_L, m^2_{\tilde{t}'} + m^2_{\tilde{X}})$$

$$C_{1,2} = \lambda_2^*(Z^i_{\tilde{t}'}^i, m^2_{\tilde{X}}, m^2_{\tilde{t}'} + m^2_{\tilde{X}})$$

(21)

and the relevant coefficients are

$$a_3 = \lambda_2^*(Z^i_{\tilde{t}'}^i), \quad a_4 = \lambda_1^*(Z^i_{\tilde{t}'}^i),$$

$$b_3 = \lambda_2^*(Z^i_{\tilde{t}'}^i), \quad b_4 = \lambda_2^*(Z^i_{\tilde{t}'}^i).$$

(22)

In Fig. 2 we present the relevant self-energy diagrams of the rare decay $t \to ch$ in BLMSSM.

As in Ref. [8], it is convenient to define the following structure:

$$\Sigma_{tc}(k) \equiv k\Sigma_L (k^2) P_L + k\Sigma_R (k^2) P_R + m_t (\Sigma_{Ls}(k^2) P_L + \Sigma_{Rs}(k^2) P_R).$$

(23)

Here $m_t$ factor is inserted there only to preserve the same dimensionality for the different $\Sigma[8]$. And the effective interaction vertex of the mixed self-energy diagrams could be taken on the following general form in terms of the various $\Sigma$. 

10
\[-iT_{Sc} = \frac{-i\gamma_{m_t}}{2m_W \sin \beta m_c^2 - m_t^2} c(p) \{ \]

\[
(P_L \cos \alpha [m_c^2 \Sigma_R(m_t^2) + m_t m_t (\Sigma_{Rs}(m_t^2) + \Sigma_L(m_c^2)) + m_t^2 \Sigma_{Ls}(m_c^2)]
+ P_R \cos \alpha [L \leftrightarrow R]) \}
\]

\[-iT_{St} = \frac{-i\gamma_{m_t}}{2m_W \sin \beta m_c^2 - m_t^2} c(p) \{ \]

\[
(P_L \cos \alpha [m_t (\Sigma_L(m_c^2) + \Sigma_{Rs}(m_t^2)) + m_c (\Sigma_R(m_c^2) + \Sigma_{Ls}(m_t^2))]
+ P_R \cos \alpha [L \leftrightarrow R]) \} t(p') \]

Comparing with Eq. 16, the corresponding contribution to the form factors $F_L$ and $F_R$ is transparent.

Using the couplings above, we could get the $\Sigma$ of self-energy diagrams in Fig.2(a) is

\[
\Sigma_L(k^2) = \frac{i}{16\pi^2} \sum_{i, l} a_1 b_2 (B_0(k^2, m_{X_l}^2, m_{l'}) + B_1(k^2, m_{X_l}^2, m_{l'}))
\]

\[
\Sigma_R(k^2) = \frac{i}{16\pi^2} \sum_{i, l} a_2 b_1 (B_0(k^2, m_{X_l}^2, m_{l'}) + B_1(k^2, m_{X_l}^2, m_{l'}))
\]

\[
m_t \Sigma_{Ls}(k^2) = \frac{i}{16\pi^2} \sum_{i, l} a_2 b_2 m_{l'} B_0(k^2, m_{X_l}^2, m_{l'})
\]

\[
m_t \Sigma_{Rs}(k^2) = \frac{i}{16\pi^2} \sum_{i, l} a_1 b_1 m_{l'} B_0(k^2, m_{X_l}^2, m_{l'}) \tag{25}
\]

with $B_{0,1}$ are the two-point functions. Similarly, the $\Sigma$ of self-energy diagrams in Fig.2(b) have the form:

\[
\Sigma_L(k^2) = \frac{i}{16\pi^2} \sum_{i} a_3 b_4 (B_0(k^2, m_{l'}^2, m_{X_l}^2) + B_1(k^2, m_{l'}^2, m_{X_l}^2))
\]

\[
\Sigma_R(k^2) = \frac{i}{16\pi^2} \sum_{i} a_4 b_3 (B_0(k^2, m_{l'}^2, m_{X_l}^2) + B_1(k^2, m_{l'}^2, m_{X_l}^2))
\]

\[
m_t \Sigma_{Ls}(k^2) = \frac{i}{16\pi^2} \sum_{i} a_4 b_4 m_{l'} B_0(k^2, m_{l'}^2, m_{X_l}^2)
\]

\[
m_t \Sigma_{Rs}(k^2) = \frac{i}{16\pi^2} \sum_{i} a_3 b_3 m_{l'} B_0(k^2, m_{l'}^2, m_{X_l}^2) \tag{26}
\]
IV. NUMERICAL ANALYSIS

In general case, the partial widths of $t \to ch$ process are

$$\Gamma(t \to ch) = \frac{g^2}{32\pi m_t^3} \lambda^{1/2}(m_t^2, m_h^2, m_c^2) \times \left[ (m_t^2 + m_c^2 - m_h^2) (|F_L|^2 + |F_R|^2) + 2m_t m_c (F_L F_R^* + F_R F_L^*) \right]$$

(27)

with $\lambda(x^2, y^2, z^2) = (x^2 - (y + z)^2)(x^2 - (y - z)^2)$ is the usual Källen function, and

$$F_L = F_L^{BLSSM} + F_L^{MSSM} + F_L^{SM}$$

$$F_R = F_R^{BLSSM} + F_R^{MSSM} + F_R^{SM}$$

(28)

In our calculation, we will use the form factors of MSSM $F_{L,R}^{MSSM}$ mentioned in [8]. And since the contributions of SM is too small, about $10^{-13}$ [7], so we ignore the form factors of SM.

To compute the branching ratio, we take the SM charged-current two-body decay $t \to bW$ to be the dominant $t$-quark decay mode, which is $\Gamma(t \to bW^+) = 1.466 |V_{tb}|^2$. We will then approximate the branching ratio by

$$Br(t \to ch) = \frac{\Gamma(t \to ch)}{\Gamma(t \to bW^+)}$$

(29)

To reduce the number of free parameters in our numerical analysis, the parameters are adopted as Ref. [23, 24]. In this choice, it is easy for the $2 \times 2$ CP-even Higgs mass squared matrix to predict the lightest eigenvector with a mass $125.9$ GeV, and the choice is good for the behavior of $h \to \gamma\gamma$ and $h \to VV^*$ ($V = Z, W$) [23].
FIG. 3: The branching ratio of $t \to ch$ varying with $m_{\tilde{Q}_4}$

\[ B_4 = \frac{3}{2}, \quad v_B = \sqrt{v_B^2 + \bar{v}_B^2} = 3\text{TeV}, \]
\[ \tan \beta = \tan \beta_B = 2, \]
\[ m_{\tilde{U}_4} = m_{\tilde{Q}_5} = m_{\tilde{U}_5} = 1\text{TeV}, \]
\[ A_{u_4} = A_{u_5} = 500\text{GeV}, \]
\[ A_{BU} = 1\text{TeV}, \quad \lambda_u = 0.5, \]
\[ Y_{u_4} = 0.76Y_t, \quad Y_{d_4} = 0.7Y_b, \]
\[ Y_{u_5} = 0.7Y_b, \quad Y_{d_5} = 0.13Y_t, \]
\[ \mu = -800\text{GeV} \]
\[ B_X = 500\text{GeV}, \quad \mu_X = 2\text{TeV}, \]

Choosing $m_{ZB} = 1\text{TeV}, \mu_B = 500\text{GeV}, \lambda_Q = 0.5, A_{BQ} = 1\text{TeV}$. We plot in Fig.3 the
BRs of $t \to ch$ versus $m_{\tilde{Q}_4}$, the solid line, dash line and dot line correspond to $\lambda_1 = \lambda_2 = 0.6, 0.4, 0.2$, respectively. We could see that the BRs decrease as $m_{\tilde{Q}_4}$ runs from 700GeV to 1300GeV, and increase when $\lambda_1 = \lambda_2$ increase, because $m_{\tilde{Q}_4}$ is the mass parameter of the exotic quarks, and $\lambda_1, \lambda_2$ proportional to the coupling coefficient. In addition, when $m_{\tilde{Q}_4} \geq 1100$, the BRs is tend to the results of MSSM.

In Fig. 4, we plot $\text{Br}(t \to ch)$ varying with $m_{Z_B}$. Adopting $m_{\tilde{Q}_4} = 790\text{GeV}, \mu_B = 500\text{GeV}, \lambda_Q = 0.5, A_BQ = 1\text{TeV}$, and with $\lambda_1 = \lambda_2 = 0.6$ (solid line), $\lambda_1 = \lambda_2 = 0.4$ (dash line), $\lambda_1 = \lambda_2 = 0.2$ (dot line). We could see that the BRs decrease as $m_{Z_B}$ runs from 800GeV to 1100GeV, since $m_{Z_B}$ contribute to the mass matrix of exotic squarks, and increase when $\lambda_1 = \lambda_2$ increase. And when $\lambda_1 = \lambda_2 = 0.6, 0.4$, $\text{Br}(t \to ch)$ is at the order of $10^{-4}$, when $\lambda_1 = \lambda_2 = 0.2$, $\text{Br}(t \to ch)$ is at the order of $10^{-5}$.

We assume $m_{\tilde{Q}_4} = 790\text{GeV}, m_{Z_B} = 1\text{TeV}, \lambda_Q = 0.5, A_BQ = 1\text{TeV}$. We plot in Fig.5 the BRs of $t \to ch$ versus $\mu_B$, the solid line, dash line and dot line correspond to $\lambda_1 = \lambda_2 =$
0.6, 0.4, 0.2, respectively. We could see that the BRs increase as $\mu_B$ runs from 300GeV to 600GeV, since $\mu_B$ inversely to the mass of exotic squarks.

Choosing $m_{\tilde{Q}_4} = 790$GeV, $m_{\tilde{Z}_B} = 1$TeV, $\mu_B = 500$GeV, $A_{BQ} = 1$TeV, we draw Br($t \to ch$) varying with $\lambda_Q$ in Fig.6 for $\lambda_1 = \lambda_2 = 0.6, 0.4, 0.2$ respectively. We could see that the curve first increase and then decrease, but not significantly, since $\lambda_Q$ contribute both to the mass of exotic squarks and the coupling coefficient.

Taking $m_{\tilde{Q}_4} = 790$GeV, $m_{\tilde{Z}_B} = 1$TeV, $\mu_B = 500$GeV, $\lambda_Q = 0.5$, we show the Br($t \to ch$) varying with $A_{BQ}$ in Fig.7 for $\lambda_1 = \lambda_2 = 0.6$ (solid line), $\lambda_1 = \lambda_2 = 0.4$ (dash line), $\lambda_1 = \lambda_2 = 0.2$ (dot line), respectively. We could see that the BRs decrease as $A_{BQ}$ runs from 1TeV to 1.8TeV, since $A_{BQ}$ contribute to the mass matrix of exotic squarks. And when $\lambda_1 = \lambda_2 = 0.6, 0.4$, Br($t \to ch$) is at the order of $10^{-4}$, when $\lambda_1 = \lambda_2 = 0.2$, Br($t \to ch$) is at the order of $10^{-5}$.
FIG. 6: The branching ratio of \( t \to ch \) varying with \( \lambda_Q \)

V. SUMMARY

The running LHC is a top-quark factory, and provides a great opportunity to seek out top-quark decays. And it is showed that the channel \( t \to ch \) could be detectable reaching a sensitivity level of \( \text{Br}(t \to ch) \sim 5 \times 10^{-5} \) \cite{26, 27}. But the fact is that the branching ratio of the process is so small in the SM\cite{8}, which is \( \text{Br}(t \to ch) \sim 10^{-13} \), so it is too small to be measurable in the near future.

In this work, we study the rare top decay to a 125GeV Higgs in the framework of the BLMSSM. Adopting reasonable assumptions on the parameter space, we present the radiative correction to the process in BLMSSM, and draw some curves between the BRs and new physics parameters. We get the branching ratio of \( t \to ch \) can reach \( 10^{-3} \), so this process could be detected in near future at LHC.
In addition, the author of [28] yields an estimated upper limit of $\text{Br}(t \to ch) < 2.7\%$ for a Higgs boson mass of 125GeV, by combining the CMS results from a number of exclusive three- and four-lepton search channels. And the ATLAS find the limit of $\text{Br}(t \to ch) < 0.83\%$ at 95% C.L. by searching for $t \to ch$, with $h \to \gamma\gamma$, in $tt$ events. [29, 30] And our numerical evaluations indicates the BRs is highly dependent upon the parameters $\lambda_{1,2}$, the sensitive parameters can make the contribution to $\text{Br}(t \to ch)$ sizeable. Considering the experiment upper bounds from CMS and ATLAS, the parameters $\lambda_{1,2}$ should not be too large under our assumptions of the parameter space.

As we could see above, the $t \to ch$ process may be found in near future, and further more constraints of BLMSSM can be obtained from more precise determinations.
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