Rare top quark decays in extended models

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Abstract. Flavor changing neutral currents (FCNC) decays $t \rightarrow H^0 + c$, $t \rightarrow Z + c$, and $H^0 \rightarrow t + \bar{c}$ are discussed in the context of Alternative Left-Right symmetric Models (ALRM) with extra isosinglet heavy fermions where FCNC decays may take place at tree-level and are only suppressed by the mixing between ordinary top and charm quarks, which is poorly constraint by current experimental values. The non-manifest case is also briefly discussed.

INTRODUCTION

Flavor-changing neutral currents (FCNC) are absent in the Standard Model (SM) at the tree-level due to the Glashow-Iliopoulos-Maiani (GIM) mechanism. However, new FCNC states can appear in top decays if there is physics beyond the Standard Model. In this context, rare top quark decays are interesting because they might be a source of possible new physics effects. In some particular models beyond the SM, rare top decays may be significantly enhanced to reach detectable levels [1].

Rare top decays have been studied in the context of the SM and beyond [2, 3, 4]. The top quark decays into gauge bosons ($t \rightarrow c + V; V \equiv \gamma, Z, g$) are extremely rare events in the SM. However, by considering physics beyond the SM, for example, the Minimal Supersymmetric Standard Model (MSSM) or the two-Higgs-doublet model (2HDM) or extra quark singlets, new possibilities open up [1, 5], enhancing this branching ratios to the order of $\sim 10^{-6}$ for the $t \rightarrow c + Z$ [6] channel and $\sim 10^{-4}$ for the $t \rightarrow c + H$ [7] case.

In the future CERN Large Hadron Collider (LHC), about $10^7$ top quark pairs will be produced per year [8]. An eventual signal of FCNC in the top quark decay will have to be ascribed to new physics. Furthermore, since the Higgs boson could also be produced at significant rates in future colliders, it is also important to search for all the relevant FCNC Higgs decays.

On the other hand, while the electroweak SM has been successful in the description of low-energy phenomena, it leaves many questions unanswered. One of them has to do with the understanding of the origin of parity violation in low-energy weak
interaction processes. Within the framework of left-right symmetric models, based on the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, this problem finds a natural answer [9, 10]. Moreover, new formulations of this model have been considered in which the fermion sector has been enlarged to include isosinglet vectorlike heavy fermions in order to explain the mass hierarchy [11, 12]. Most of these models includes two Higgs doublets.

We consider the rare top decay into a Higgs boson and the FCNC decay of the Higgs boson with the presence of a top quark in the final state, within the context of these alternative left-right models (ALRM) with extra isosinglet heavy fermions. Due to the presence of extra quarks the Cabibbo-Kobayashi-Maskawa matrix is not unitary and FCNC may exist at tree-level.

THE MODEL

The ALRM formulation is based on the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. In order to solve different problems such as the hierarchy of quark and lepton masses or the strong CP problem, different authors have enlarged the fermion content to be of the form

$$l^0_{iL} = \begin{pmatrix} v^0_i \\ e^0_i \end{pmatrix}_L, \quad e^0_{iR}, \quad \tilde{l}^0_{iL} = \begin{pmatrix} \tilde{v}^0_i \\ \tilde{e}^0_i \end{pmatrix}_L, \quad \tilde{e}^0_{iR},$$

$$Q^0_{iL} = \begin{pmatrix} u^0_i \\ d^0_i \end{pmatrix}_L, \quad u^0_{iR}, \quad d^0_{iR}, \quad \tilde{Q}^0_{iR} = \begin{pmatrix} \tilde{u}^0_i \\ \tilde{d}^0_i \end{pmatrix}_R, \quad \tilde{d}^0_{iL}, \quad \tilde{d}^0_{iL},$$

where the index $i$ ranges over the three fermion families. The superscript 0 denote weak eigenstates. In many of these models, extra neutral leptons also appears in order to explain the neutrino mass pattern, however we will focus in this work only on the quark sector.

In order to break $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ down to $U(1)_{em}$ the ALRM introduces two Higgs doublets, the SM one ($\phi$) and its partner ($\tilde{\phi}$). Ref. [13] shows that from the eight scalar degrees of freedom, six become the Goldstone bosons required to give mass to the $W^\pm, \tilde{W}^\pm, Z$ and $\tilde{Z}$; thus two neutral Higgs bosons, $H$ and $\tilde{H}$, remain in the physical spectrum.

The renormalizable and gauge invariant interactions of the scalar doublets $\phi$ and $\tilde{\phi}$ with the fermions are described by the Yukawa Lagrangian. For the quark fields, the corresponding Yukawa terms are written as

$$L^q_i = \lambda^d_{ij} Q^0_{il} \phi d^0_{jr} + \lambda^u_{ij} Q^0_{ir} \tilde{\phi} u^0_{jr} + \tilde{\lambda}^d_{ij} \tilde{Q}^0_{ir} \tilde{\phi} d^0_{jl} + \mu^d_{ij} \tilde{Q}^0_{ir} \phi d^0_{jr} + \mu^u_{ij} u^0_{ir} \tilde{u}^0_{jr} + h.c.$$  \hspace{1cm} (2)

where $i, j = 1, 2, 3$ and $\lambda^d_{ij}, \tilde{\lambda}^d_{ij}, \lambda^u_{ij}$, and $\mu^d_{ij}$ are (unknown) matrices. The conjugate fields $\tilde{\phi}$ ($\tilde{\phi}$) are $\tilde{\phi} = i\tau_2 \phi^*$ and $\tilde{\phi} = i\tau_2 \tilde{\phi}^*$, with $\tau_2$ the Pauli matrix.
We can introduce the generic vectors $\psi^0_L$ and $\psi^0_R$ [14], for representing left and right electroweak states with the same charge. These vectors can be decomposed into the ordinary $\psi^0_{OL}$ and the exotic $\psi^0_{EL}$ sectors

$$
\psi^0_L = \begin{pmatrix} \psi^0_{OL} \\ \psi^0_{EL} \end{pmatrix}, \quad \psi^0_R = \begin{pmatrix} \psi^0_{OR} \\ \psi^0_{ER} \end{pmatrix},
$$

(3)

In the same way we can define the vectors for the mass eigenstates in terms of ‘light’ $\psi_{fL}$ and ‘heavy’ $\psi_{hL}$ states. The relation between weak eigenstates and mass eigenstates will be given through the matrices $U_L$ and $U_R$ by $\psi^0_L = U_L \psi_L, \psi^0_R = U_R \psi_R$ where

$$
U_a = \begin{pmatrix} A_a & E_a \\ F_a & G_a \end{pmatrix}, \quad a = L, R
$$

(4)

Here, $A_a$ is the $3 \times 3$ matrix relating the ordinary weak states with the light-mass eigenstates, $G_a$ is a $3 \times 3$ matrix relating the exotic states with the heavy ones, while $E_a$ and $F_a$ describe the mixing between the two sectors.

In this model, thanks to the extra heavy quarks, it is possible to have a relatively big mixing between ordinary quarks. This is not a particular characteristic of the model but a general feature when considering models with extra heavy singlets [15].

The tree-level interaction of the neutral Higgs bosons $H$ and $\tilde{H}$ with the light fermions are given by

$$
\mathcal{L}_f^f = \frac{g}{2\sqrt{2}} \bar{\psi} A^\dagger_f A_f \frac{m_f}{M_W} \psi R \left( H \cos \alpha - \tilde{H} \sin \alpha \right) + \frac{\tilde{g}}{\sqrt{2}} \bar{\psi} F^\dagger_f F_f \psi R \left( H \sin \alpha + \tilde{H} \cos \alpha \right) + h.c.
$$

(5)

The neutral current in terms of the mass eigenstates, including the contribution of the neutral gauge boson mixing, can be written directly from this Lagrangian.

From the last equation we can see that, thanks to the non-unitarity of the $A_a$ matrices we can have FCNC at tree-level. This characteristic appears due to the extra quark content of the model, which is not present in the usual left-right symmetric model.

**FCNC TOP AND HIGGS DECAYS IN THE ALRM**

Once we have introduced the model in which we are interested, we compute the expected branching ratio for a FCNC top or Higgs decay with a charm quark in the final state. We perform this analysis in this section.

**Constraining the top-charm mixing angle**

In order to have an expectation on the branching ratio for the FCNC top decay in the ALRM we need first an estimate on the mixing between the top and charm quarks in
the model. One may think that the best constrain could come from the flavor-changing coupling of the neutral Z boson to the top and charm quarks, which can be written as:

\[ \mathcal{L}_Z^{ct} = \frac{e}{s_{\theta_w} c_{\theta_w}} (g_V + g_A \gamma^5) \gamma^\mu Z \mu t \]  

(6)

where

\[ g_{V,A} = \frac{1}{4} \left( c_{\Theta} - \frac{g_{\Theta}^2}{g r_{\theta_w}} s_{\Theta} \right) \eta_{32}^L \pm \frac{1}{4} \frac{g_{\Theta}^2}{g r_{\theta_w}} s_{\Theta} \eta_{32}^R \]  

(7)

and \( s_{\theta_w}, c_{\theta_w} \) and \( r_{\theta_w} \) are, respectively, \( \sin \theta_w, \cos \theta_w \) and \( \sqrt{\cos^2 \theta_w - (g_{\Theta}^2/g)^2 \sin^2 \theta_w} \); \( \theta_w \) is the weak mixing angle, \( \Theta \) is the mixing between the \( Z \) and \( Z' \) neutral gauge bosons. Here, \( \eta_{32}^L \) and \( \eta_{32}^R \) represent the mixing between the ordinary top and charm quarks and are given by

\[ \eta_{32}^L = (A_L^L A_L)_{32} \quad \eta_{32}^R = (A_R^L A_R)_{32}. \]  

(8)

The mixing between the \( Z \) and the \( Z' \) neutral gauge bosons, \( \Theta \), is expected to be small [16] if the ratio \( r_g = g/\bar{g} = 1 \). However, one might think that for different values of \( r_g \) these bounds are not longer valid. This is actually true, however, for most of the values of \( r_g \), the expected freedom for the mixing angle \( \Theta \) is still limited.

We can see from the definition of \( r_{\theta_w} \) that the value of this ratio can not be bigger than \( \sqrt{\cos^2 \theta_w/\sin^2 \theta_w} \approx 1.82 \). We can recalculate the constraint obtained in [16] taking into account the freedom of this parameter. In order to do this analysis we simply need to consider the appropriate range for the parameter \( r_g \) that will affect the coupling constants for the \( Z \to e^+e^- \) that are needed for such computation:

\[ g_{V,A}^e = -\frac{1}{2} \left( c_{\Theta} - \frac{g_{\Theta}^2}{g r_{\theta_w}} s_{\Theta} \right) \pm \frac{1}{2} \frac{g_{\Theta}^2}{g r_{\theta_w}} s_{\Theta}. \]  

(9)

Note that in this case we are not taking into account the lepton flavor violation that has been discussed in ref [13].

With this formula and the limit for \( g_A \) obtained from the experiment: \( g_A^{\exp} = -0.4998 \pm 0.00014 \) [17] we can obtain a constraint for the \( \Theta \) mixing depending on the value of \( r_g \). The result of such analysis is shown in Fig. (1) were it is possible to see that \( |\sin \Theta| \leq 0.03 \), if the value of the ratio \( r_g \) is smaller than 1.6. As \( r_g \) approaches the critical value of 1.8, the constraint will disappear.

Therefore, the mixing angle \( \Theta \) can be safely neglected for most of the values of \( g \) and \( \bar{g} \). In this case the expression in Eq. (7) will not depend on the parameter \( \eta_{32}^R \). For the present analysis we will consider only the case with \( g = \bar{g} \), and therefore from now on we will denote \( \eta_{32} = \eta_{32}^L \).

From Eq. (6) we can compute the branching ratio for the decay \( t \to Z + c \) and compare it to the experimental limit \( B(t \to Z + c) \leq 0.137 \) [18] at 95 % C. L. We will get the maximum value for \( \eta_{32} \leq 0.53 \).

Although we have found a direct constrain to \( \eta_{32} \), it is possible to get a stronger limit if we use the unitarity properties of the mixing matrix and the constrain on \( \eta_{22} \) that comes from the branching ratio \( \Gamma(Z \to c + \bar{c}) \). The experimental value for the branching
FIGURE 1. Expected value of $g_A$ for the model in dependence on the mixing angle between $Z$ and $\hat{Z}$, and the ratio $r_g = g/\hat{g}$. The horizontal lines shows the 90 % C. L. allowed by the experiment.

ratio of this process is given by $B(Z \to c\bar{c}) = \Gamma(Z \to c\bar{c})/\Gamma_{total} = 0.1181 \pm 0.0033$ (see [19]). Using this experimental value, the minimum value for $\eta_{22}$ at 95 % C. L. will be $\eta_{22} \geq 0.99$.

This information is of great help for constraining $\eta_{32}$ since the unitarity of the mixing matrix has already been analyzed in the general case [20] and leads to the relation $|\eta_{32}|^2 \leq (1 - \eta_{33})(1 - \eta_{22})$.

Although we don’t know the value for $\eta_{33}$, the boundary on $\eta_{22}$ is enough to see that the mixing parameter $\eta_{32} \leq 0.1$. The higher value $\eta_{23} = 0.1$ is obtained when we take the extreme case $\eta_{33} = 0$, as can be seen from the equation in the previous paragraph.

It is possible to obtain more stringent constraints if low-energy data are considered. For the case of two extra quark singlets, this analysis was done in a very general framework in Ref. [5]. After a very complete analysis of all the observables, the author of this article obtained $|\eta_{32}| \leq 0.036$. This relatively large value is allowed for the case of a exotic top mass similar to that of the SM top-quark (There are not stringent lower bounds on the mass of a exotic top quark, being 220 GeV the current direct limit [21]). In the case of a very heavy mass for the exotic top-quark the constraint is more stringent: $|\eta_{32}| \leq 0.009$. In what follows we will use these two values in order to illustrate the expected signals from rare Higgs and top decays.

The decays $t \to H^0 + c$ and $H^0 \to t + \bar{c}$

Now that we have an estimate for the value of $\eta_{32}$, we compute the branching ratio for $t \to H^0 + c$ in the framework of ALRM. We take the charged-current two-body decay $t \to b + W$ to be the dominant $t$-quark decay mode. The neutral Higgs boson $H^0$ will be assumed to be the lightest neutral mass eigenstate.
Assuming $M_{H^0} \gg M_H$ the vertex $t c H^0$ is written as $\frac{g m_{H^0} \eta_{32}}{2 m_W} \cos \alpha P_L$. The partial width for this tree-level process can be obtained in the usual way and it is given by:

$$
\frac{G_F \eta_{32} \cos^2 \alpha}{16 \sqrt{2} \pi m_t} \left( m_t^2 + m_c^2 - M_{H^0}^2 \right) \left[ \left( m_t^2 - \left( M_{H^0} + m_c \right)^2 \right) \left( m_t^2 - \left( M_{H^0} - m_c \right)^2 \right) \right]^{\frac{1}{2}}
$$

where $G_F$ is the Fermi’s constant, $m_t$ denotes the top mass, $m_c$ is the charm mass, and $M_{H^0}$ is the mass of the neutral Higgs boson. We can see from this formula that the branching ratio will be proportional to the product $\eta_{32} \cos \alpha$, of the top-quark mixing with the SM Higgs boson mixing with the extra Higgs boson.

The branching ratio for this decay is obtained as the ratio of Eq. (10) to the total width for the top quark, namely $B(t \to H^0 + c) = \frac{\Gamma(t \to H^0 + c)}{\Gamma(t \to b + W)}$.

Thanks to the possible combined effect of a big $\cos \alpha$ (null mixing between the SM Higgs boson and the additional Higgs bosons) and a big value of $\eta_{32}$ this branching ratio could be as high as $3 \times 10^{-4}$, for a Higgs mass of 117 GeV. Perhaps is more realistic to consider the more stringent constraint $\eta_{32} = 0.009$, but even in this case, for $\cos \alpha \approx 1$ there is still sensitivity for detecting a positive signal of order $10^{-5}$.

Finally we also consider the case of a Standard Higgs with a large mass. The best-fit value of the expected Higgs mass, including the new average for the mass of the top quark, is 117 GeV [22] and the upper bound is $M_{H^0} \leq 251$ GeV at 95 % C. L. However, the error for the Higgs boson mass from this global fit is asymmetric, and a Higgs mass of 400 GeV is well inside the 3σ region as can be seen in Ref [22].

We estimate the branching ratio for the decay $H^0 \to t + \bar{c}$, where $H^0$ is the light neutral Higgs boson of the ALRM. The expression for the partial width is

$$
\frac{3 G_F m_t^2 \eta_{32}^2 \cos^2 \alpha}{8 \sqrt{2} \pi M_{H^0}^2} \left( M_{H^0}^2 - m_t^2 - m_c^2 \right) \left[ \left( M_{H^0}^2 - \left( m_t + m_c \right)^2 \right) \left( M_{H^0}^2 - \left( m_c - m_t \right)^2 \right) \right]^{\frac{1}{2}}.
$$

The branching ratio for this decay is obtained as the ratio of Eq. (11) to the total width of the Higgs boson, which will include the dominant modes $H^0 \to b + \bar{b}$, $H^0 \to c + \bar{c}$, $H^0 \to t + \bar{t}$, $H^0 \to W + W$, and $H^0 \to Z + Z$. The expressions for these decay widths in the ALRM also includes corrections due to the new parameters introduced in the model, and they are taken into account [23].

We computed the branching ratios for different decay modes, both for the Standard Model case ($\eta_{32} = 0$ and $\eta_{ii} = 1$) and for the FCNC case. We found that, also for a heavy Higgs, there are chances to either detect or to constrain the mixing angle parameter $\eta_{32}$. In this case, since all the partial widths have the same dependence on $\cos^2 \alpha$, the branching ratios will depend only on $\eta_{32}$.

**RESULTS AND CONCLUSIONS**

The ALRM allows relatively big values of $\eta_{32}$. The $t \to H + c$ branching ratio could be of order of $10^{-4}$, which is at the reach of LHC. It has been estimated that the LHC sensitivity (at 95 % C. L.) for this decay is $Br(t \to Hc) \leq 4.5 \times 10^{-5}$ [24]; this branching
ratio would be obtained in this model for a top-charm mixing $\eta_{32} = 0.015$ and a diagonal ordinary top coupling $\eta_{22} \simeq 0.98$. On the other hand, the FCNC mode $H \rightarrow t + \bar{c}$ may reach a branching ratio of order $10^{-3}$ and can also be a useful channel to look for signals of physics beyond the SM in the LHC.

ACKNOWLEDGMENTS

This work has been supported by Conacyt and SNI.

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