Rare top quark and Higgs boson decays in alternative left-right symmetric models

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Top quark and Higgs boson decays induced by flavor-changing neutral currents (FCNC) are very much suppressed in the standard model. Their detection in colliders such as the Large Hadron Collider, Next Linear Collider, or Tevatron would be a signal of new physics. We evaluate the FCNC decays $t \to H^0 + c$, $t \to Z + c$, and $H^0 \to t + \bar{c}$ in the context of alternative left-right symmetric models with extra isosinglet heavy fermions; in this case, FCNC decays occur at tree level, and they are suppressed only by the mixing between ordinary top and charm quarks, which is poorly constrained by current experimental values. This provides the possibility for future colliders either to detect new physics or to improve present bounds on the parameters of the model.

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I. INTRODUCTION

Rare top quark decays are interesting because they might be a source of possible new physics effects. Because of its large mass of about 178 GeV [1], the top quark dominant decay mode is into the channel $t \to b + W$. In the standard model (SM), based on the spontaneously broken local symmetry $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, flavor-changing neutral currents (FCNC) are absent at the tree level due to the Glashow-Iliopoulos-Maiani mechanism, and they are extremely small at loop level. However, new FCNC states can appear in top decays if there is physics beyond the standard model. Moreover, in some particular models beyond the SM, rare top decays may be significantly enhanced to reach detectable levels [2].

Rare top decays have been studied in the context of the SM and beyond [3–5]. The top quark decays into gauge bosons ($t \to c + V$; $V \equiv \gamma, Z, g$) are extremely rare events in the SM; their branching ratios are, according to Refs. [3,6], 5×10^{-13} for the photon, 10^{-13} for the Z boson, and $\sim 4 \times 10^{-11}$ for the gluon channel, and even smaller according to other estimates [7]. Similarly, the top quark decay into the SM Higgs boson is a very rare decay, with $\mathrm{BR}(t \to c + H) \sim 10^{-14}$ [5,8]. 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 [2,4–11], enhancing this branching ratio to the order of $\sim 10^{-6}$ for the $t \to c + Z$ [7] channel and $\sim 10^{-4}$ for the $t \to c + H$ [9] case. The rare top decay $t \to q + W + Z$ has also been considered as a future test of new physics [12].

On the other hand, the FCNC decays of the Higgs boson can be important in various scenarios, including the MSSM [13]. The FCNC Higgs decay into a top quark within a general 2HDM has been studied in Ref. [14]. Because the FCNC Higgs decays in the SM are very suppressed, any experimental signature of Higgs FCNC type could be evidence of physics beyond the SM.

In the future CERN Large Hadron Collider (LHC), about 10^7 top quark pairs will be produced per year [15]. 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 [16,17]. 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 [18,19], the smallness of the neutrino mass [20], or the problem of weak and strong CP violation [21,22]. Most of these models include two Higgs doublets.

In this paper, 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. Because of the presence of extra quarks, the Cabibbo-Kobayashi-Maskawa matrix is not unitary and FCNC may exist at tree level.

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Therefore, a high branching ratio for the decay $t \to c + H$ (for a Higgs boson lighter than the top quark mass) or for the decay $H \to t + \bar{c}$ is allowed, opening great opportunities either to detect or to constrain the mixing parameter η_{32} between the ordinary top and charm quarks.

The organization of the paper is as follows: In Sec. II we review the alternative left-right model, giving emphasis to the fermion mixing and flavor violation. In Sec. III we present our calculations in the ALRM for the processes $t \to c + Z$, $t \to H^0 + c$, and $H^0 \to t + \bar{c}$; we derive bounds on the parameters of the model associated with FCNC transitions and we discuss future perspectives for improving these bounds. Section IV contains our conclusions.

II. 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

$$\begin{aligned} l^0_{iL} &= \begin{pmatrix} \nu^0_i \\ e^0_i \end{pmatrix}_L, e^0_{iR}; \qquad \hat{l}^0_{iR} = \begin{pmatrix} \hat{\nu}^0_i \\ \hat{e}^0_i \end{pmatrix}_R, \hat{e}^0_{iL} \\ Q^0_{iL} &= \begin{pmatrix} u^0_i \\ d^0_i \end{pmatrix}_L, u^0_{iR}, d^0_{iR}; \qquad \hat{Q}^0_{iR} = \begin{pmatrix} \hat{u}^0_i \\ \hat{d}^0_i \end{pmatrix}_R, \hat{u}^0_{iL}, \hat{d}^0_{iL}, \end{aligned} \tag{1}$$

where the index i ranges over the three fermion families. The superscript 0 denotes weak eigenstates. The quantum numbers of these fermions, under the gauge group $SU(3)_C \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, are given by

$$\begin{aligned} &l^0_{iL} \quad (1,2,1)_{-1} \quad e^0_{iR} \quad (1,1,1)_{-2}; \\ &\hat{l}^0_{iR} \quad (1,1,2)_{-1} \quad \hat{e}^0_{iL} \quad (1,1,1)_{-2} \qquad u^0_{iR} \quad (3,1,1)_{4/3} \\ &d^0_{iR} \quad (3,1,1)_{2/3}; \qquad \hat{u}^0_{iL} \quad (3,1,1)_{4/3} \qquad \hat{d}^0_{iL} \quad (3,1,1)_{2/3} \\ &Q^0_{iL} \quad (3,2,1)_{1/3} \qquad \hat{Q}^0_{iR} \quad (3,1,2)_{1/3}. \end{aligned} \tag{2}$$

In many of these models, extra neutral leptons also appear 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 (ϕ) and its partner ($\hat{\phi}$). The symmetry breaking is done in such a way that the vacuum expectation values of the Higgs fields are

$$\langle\phi\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v \end{pmatrix}; \qquad \langle\hat{\phi}\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \hat{v} \end{pmatrix}. \tag{3}$$

Reference [23] shows that, from the eight scalar degrees of freedom, six become the Goldstone bosons required to give mass to the $W^\pm$, $\hat{W}^\pm$, Z, and $\hat{Z}$; thus, two neutral Higgs bosons remain in the physical spectrum. The neutral physical states are

$$H = \sqrt{2}(\Re e\phi^0 - v)\cos\alpha + (\Re e\hat{\phi}^0 - \hat{v})\sin\alpha, \tag{4}$$

$$\hat{H} = -\sqrt{2}(\Re e\phi^0 - v)\sin\alpha + (\Re e\hat{\phi}^0 - \hat{v})\cos\alpha, \tag{5}$$

where α denotes the neutral Higgs mixing angle (which diagonalizes the neutral Higgs mass matrix). The renormalizable and gauge invariant interactions of the scalar doublets ϕ and $\hat{\phi}$ with the fermions are described by the Yukawa Lagrangian. For the quark fields, the corresponding Yukawa terms are written as

$$\begin{aligned} \mathcal{L}^q_Y = {}& \lambda^d_{ij}\overline{Q^0_{iL}}\phi d^0_{jR} + \lambda^u_{ij}\overline{Q^0_{iL}}\tilde{\phi} u^0_{jR} + \hat{\lambda}^d_{ij}\overline{\hat{Q}^0_{iR}}\hat{\phi}\hat{d}^0_{jL} \\ &+ \hat{\lambda}^u_{ij}\overline{\hat{Q}^0_{iR}}\tilde{\hat{\phi}}\hat{u}^0_{jL} + \mu^d_{ij}\overline{\hat{d}^0_{iL}}d^0_{jR} + \mu^u_{ij}\overline{\hat{u}^0_{iL}}u^0_{jR} + \text{H.c.}, \end{aligned} \tag{6}$$

where $i, j = 1, 2, 3$ and $\lambda^{d(u)}_{ij}$, $\hat{\lambda}^{d(u)}_{ij}$, and $\mu^{d(u)}_{ij}$ are (unknown) matrices. The conjugate fields $\tilde{\phi}$ ($\tilde{\hat{\phi}}$) are $\tilde{\phi} = i\tau_2\phi^*$ and $\tilde{\hat{\phi}} = i\tau_2\hat{\phi}^*$, with τ_2 the Pauli matrix.

We can introduce the generic vectors [24] ψ^0_L and ψ^0_R for representing left and right electroweak states with the same charge. These vectors can be decomposed into the ordinary and exotic sector by

$$\psi^0_L = \begin{pmatrix} \psi^0_{OL} \\ \psi^0_{EL} \end{pmatrix}, \qquad \psi^0_R = \begin{pmatrix} \psi^0_{OR} \\ \psi^0_{ER} \end{pmatrix}, \tag{7}$$

where ψ^0_{OL} is a column vector consisting of the SM $SU(2)_L$ doublets (for example, the u^0_{iL}), while ψ^0_{EL} contains the exotic singlets ($\hat{u}^0_{iL}$). The vector ψ^0_{OR} contains the SM singlets (such as u^0_{iR}) and ψ^0_{ER} contains the exotic $SU(2)_R$ doublets ($\hat{u}^0_{iR}$).

In the same way, we can define the vectors for the mass eigenstates in terms of "light" and "heavy" states:

$$\psi_L = \begin{pmatrix} \psi_{lL} \\ \psi_{hL} \end{pmatrix}, \qquad \psi_R = \begin{pmatrix} \psi_{lR} \\ \psi_{hR} \end{pmatrix}. \tag{8}$$

The relation between weak eigenstates and mass eigenstates will be given through the matrices U_L and U_R:

$$\psi^0_L = U_L\psi_L, \qquad \psi^0_R = U_R\psi_R, \tag{9}$$

where

$$\mathsf{U}_a = \begin{pmatrix} \mathsf{A}_a & \mathsf{E}_a \\ \mathsf{F}_a & \mathsf{G}_a \end{pmatrix}, \qquad a = L, R. \tag{10}$$

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

It is easy to see that in this case, the A_a is not necessarily unitary. Instead, the unitarity of the U_a matrices leads to the relations

$$A^\dagger_a A_a + F^\dagger_a F_a = I, \qquad A^\dagger_a A_a + E^\dagger_a E_a = I. \tag{11}$$

Therefore, 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 [25].

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

$$\mathcal{L}_Y^f = \frac{g}{2\sqrt{2}} \overline{\psi}_L A_L^\dagger A_L \frac{m_f}{M_W} \psi_R (H \cos\alpha - \hat{H} \sin\alpha) + \frac{\hat{g}}{\sqrt{2}} \overline{\psi}_L \frac{m_f}{M_{\hat{W}}} F_R^\dagger F_R \psi_R (H \sin\alpha + \hat{H}\cos\alpha) + \text{H.c.} \tag{12}$$

The neutral current in terms of the mass eigenstates, including the contribution of the neutral gauge boson mixing, can be written as follows:

$$-\mathcal{L}^{\text{n.c.}} = \sum_{a=LR} \overline{\psi}_a \gamma^\mu \mathsf{U}_a^\dagger \left(g\mathsf{T}_{3a}, \hat{g}\hat{\mathsf{T}}_{3a}, g' \frac{Y_a}{2} \right) \mathsf{U}_a \psi_a \begin{pmatrix} Z \\ \hat{Z} \\ A \end{pmatrix}, \tag{13}$$

where T_{3a}, $\hat{\mathsf{T}}_{3a}$, and Y are the generators of the $SU(2)_L$, $SU(2)_R$, and $U(1)_{B-L}$, respectively. For the sake of simplicity, we will consider only the case $g = \hat{g}$.

From the last two equations, we can see that, thanks to the nonunitarity 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.

III. 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. We will start by searching the maximum allowed value for a top-charm mixing and then we will obtain the possible branching ratio both for the top decay into a Higgs boson plus a charm quark and for the Higgs decay into a top plus an anticharm quark.

A. 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 the charm quarks in the model. One may think that the best constraint could come from the flavor-changing coupling of the neutral Z boson to the top and the charm quarks, which can be written as:

$$\mathcal{L}_Z^{ct} = \frac{e}{s_{\theta_W} c_{\theta_W}} \overline{c}(g_V + g_A)\gamma^\mu Z_\mu t, \tag{14}$$

where

$$g_{V,A} = \frac{1}{4}\left(c_\Theta - \frac{s_{\theta_W}^2}{r_{\theta_W}} s_\Theta\right)\eta_{32}^L \pm \frac{1}{4}\frac{c_{\theta_W}^2}{r_{\theta_W}} s_\Theta \eta_{32}^R, \tag{15}$$

and s_{θ_W}, c_{θ_W}, and r_{θ_W} are, respectively, $\sin\theta_W$, $\cos\theta_W$, and $\sqrt{\cos^2\theta_W - \sin^2\theta_W}$; θ_W is the weak mixing angle, Θ is the mixing between the Z and $\hat{Z}$ neutral gauge bosons. Here η_{32}^L and η_{32}^R represent the mixing between the ordinary top and charm quarks and are given by

$$\eta_{32}^L = (A_L^+ A_L)_{32}, \qquad \eta_{32}^R = (A_R^+ A_R)_{32}. \tag{16}$$

Since the mixing between the Z and the $\hat{Z}$ neutral gauge bosons, Θ, is expected to be small [26], it can be safely neglected, and this partial width will not depend on the parameter η_{32}^R. Therefore, from now on we will denote $\eta_{32} = \eta_{32}^L$.

From Eq. (14) we can compute the branching ratio for the decay $t \to Z + c$ and compare it to the experimental limit $B(t \to Z + c) \le 0.137$ [27] at 95% C.L. We will get the maximum value for $\eta_{32} \le 0.53$.

Although we have found a direct constraint to η_{32}, it is possible to get a stronger limit if we use the unitarity properties of the mixing matrix and the constraint on η_{22} that comes from the branching ratio $\Gamma(Z \to c + \bar{c})$. The experimental value for the branching ratio of this process is given by $B(Z \to c\bar{c}) = \Gamma(Z \to c\bar{c})/\Gamma_{\text{total}} = 0.1181 \pm 0.0033$ (see [28]). Using this experimental value, the minimum value for η_{22} at 95% C.L. will be $\eta_{22} \ge 0.99$.

This information is of great help for constraining η_{32}, since the unitarity of the mixing matrix has already been analyzed in the general case [29] and leads to the following relation:

$$|\eta_{32}|^2 \le (1 - \eta_{33})(1 - \eta_{22}). \tag{17}$$

Although we do not know the value for η_{33}, the boundary on η_{22} is enough to see that the mixing parameter $\eta_{32} \le 0.1$. The higher value $\eta_{23} = 0.1$ is obtained when we take the extreme case $\eta_{33} = 0$, as can be seen from Eq. (17).

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. [11]. After a very complete analysis of all the observables, the author of this article obtained $|\eta_{32}| \le 0.036$. This relatively large value is allowed for the case of an exotic top mass similar to that of the SM top quark [30]. In the case of a very heavy mass for the exotic top quark, the constraint is more stringent: $|\eta_{32}| \le 0.009$. In what follows, we will use these two values in order to illustrate the expected signals from rare Higgs and top decays.

B. The decay $t \to H^0 + c$

Now that we have an estimate for the value of η_{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 $M_{\hat{H}} \gg M_H$, the vertex tcH^0 is written as follows:

$$\frac{g m_t \eta_{32}}{2M_W} \cos\alpha P_L. \tag{18}$$

The partial width for this tree-level process can be obtained in the usual way and it is given by:

$$\Gamma(t \to H^0 + c) = \frac{G_F \eta_{32}^2 \cos^2\alpha}{16\sqrt{2}\pi m_t}(m_t^2 + m_c^2 - M_H^2)[(m_t^2 - (M_H + m_c)^2)(m_t^2 - (M_H - m_c)^2)]^{1/2}, \tag{19}$$

where G_F is the Fermi's constant, m_t denotes the top mass, m_c is the charm mass, and M_H 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. (19) to the total width for the top quark, namely,

$$B(t \to H^0 + c) = \frac{\Gamma(t\to H^0 + c)}{\Gamma(t \to b + W)}. \tag{20}$$

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 η_{32}, this branching ratio could be as high as $\approx 3 \times 10^{-4}$ for a Higgs mass of 117 GeV as is illustrated in Fig. 1. Perhaps it 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} as is shown in Fig. 1.

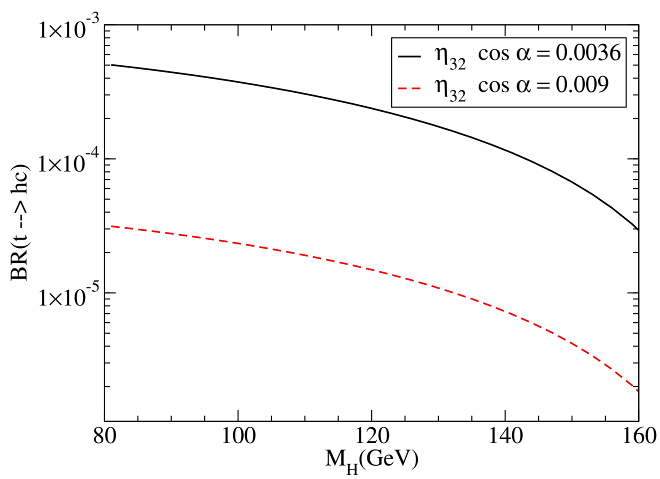

FIG. 1 (color online). Branching ratio for the rare top decay $t \to H + c$ for different values of the product of the mixing, α, between the lightest Higgs bosons and the additional Higgs boson of the model, and the mixing between the top quark and the charm quark, η_{32}. This figure shows that there is a lot of room for future collider experiments either to detect or to set bounds on these parameters.

C. The decay $H^0 \to t + \bar{c}$

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 [1] and the upper bound is $M_H \leq 251$ 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. [1].

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

$$\Gamma(H^0 \to t + \bar{c}) = \frac{3G_F m_t^2 \eta_{32}^2 \cos^2\alpha}{8\sqrt{2}\pi M_H^3}(M_H^2 - m_t^2 - m_c^2)[(M_H^2 - (m_t + m_c)^2)(M_H^2 - (m_c - m_t)^2)]^{1/2}. \tag{21}$$

The branching ratio for this decay is obtained as the ratio of Eq. (21) to the total width of the Higgs boson, which will include the dominant modes $H^0 \to b + \bar{b}$, $H^0 \to c + \bar{c}$, $H^0 \to \tau + \bar{\tau}$, $H^0 \to W + W$, and $H^0 \to Z + Z$. The expressions for these decay widths in the ALRM are

$$\Gamma(H^0 \to f + \bar{f}) = C_f \frac{G_F m_f^2 M_H \eta_{ff}^2 \cos^2\alpha}{4\sqrt{2}\pi}(1 - 4\lambda_f)^{3/2}, \tag{22}$$

where $\lambda_f = (m_f/M_H)^2$, $C_f = 1$ for leptons and $C_f = 3$ for quarks,

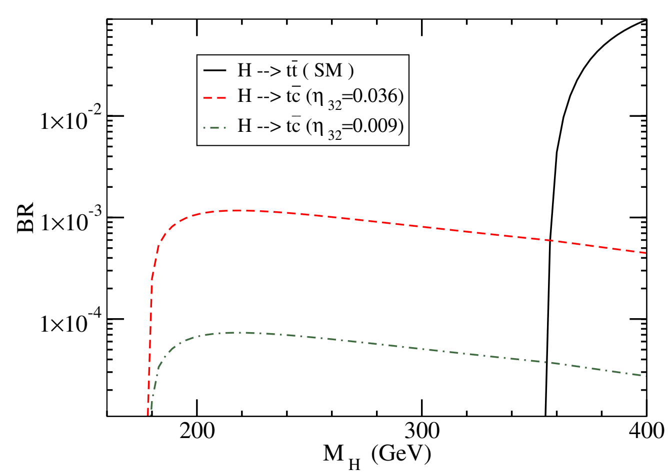

FIG. 2 (color online). Branching ratio for the rare Higgs decay $H \to t + \bar{c}$, for different values of η_{32} as a function of the Higgs mass. The standard Higgs decay $H \to t + \bar{t}$ is also shown.

$$\Gamma(H^0 \to W + W) = \frac{G_F M_H^3 \cos^2\alpha}{8\sqrt{2}\pi}(1 - 4\lambda_W)^{1/2} \times (1 - 4\lambda_W + 12\lambda_W^2), \quad (23)$$

with $\lambda_W = (M_W/M_H)^2$, and

$$\Gamma(H^0 \to Z + Z) = \frac{G_F M_H^3 M_W^4 X^2 \cos^2\alpha}{16\sqrt{2}\pi M_Z^4}(1 - 4\lambda_Z)^{1/2} \times (1 - 4\lambda_Z + 12\lambda_Z^2), \quad (24)$$

with $\lambda_Z = (M_Z/M_H)^2$ and $X = (c_{\theta_W} c_\Theta - (s_{\theta_W}/r_{\theta_W}) t_{\theta_W}(s_\Theta - r_{\theta_W} c_\Theta))^2$.

We show in Fig. 2 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 can see that, also for a heavy Higgs, there are chances to either detect or to constrain the mixing angle parameter η_{32}. In this case, since all the partial widths have the same dependence on $\cos^2\alpha$, the branching ratios will depend only on η_{32}.

IV. RESULTS AND CONCLUSIONS

We have seen that the ALRM allows relatively big values of η_{32}. The $t \to H + c$ branching ratio could be of order of 10^{-4}, which is at the reach of LHC. For example, it has been estimated that the LHC sensitivity (at 95% C.L.) for this decay is $\mathrm{Br}(t \to Hc) \leq 4.5 \times 10^{-5}$ [32]; 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 $H \to t + \bar{c}$ 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.

The ALRM is a well motivated model that rises from different theoretical motivations and has a rich phenomenology. In particular, we have studied the ALRM in the context of rare top decays, and we have found that these models could be tested in the next generation of colliders.

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