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Decay $t \to c\gamma$ in models with $SU_L(3) \times U_X(1)$ gauge symmetry

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The one-loop level mediated $t \to c\gamma$ decay is analyzed in the framework of 331 models, which are based on the $SU_L(3) \times U_X(1)$ gauge symmetry and require that the quark families transform differently in order to cancel anomalies, thereby inducing three-level flavor-changing neutral currents mediated by an extra neutral gauge boson Z' and a neutral scalar boson ϕ. These models also predict new charged gauge and scalar bosons, together with three new quarks, which can be exotic (with electric charges of $-4/3e$ and $5/3e$) or standard-model-like. Apart from the contribution of the W boson, the $t \to c\gamma$ decay receives contributions induced by the extra gauge boson and the neutral scalar boson, which are generic for 331 models. In the so-called minimal 331 model, there are additional contributions from the new charged gauge and scalar bosons accompanied by the exotic quarks. We present analytical results for the most general $t \to c\gamma$ amplitude in terms of transcendental functions. For the numerical analysis we focus on the minimal 331 model: the current bounds on the model parameters are examined and a particular scenario is discussed in which the corresponding branching ratio could be of the order of 10^{-6}, with the dominant contributions arising from the charged gauge bosons and a relatively light neutral scalar boson with flavor-changing couplings, whereas the Z' contribution would be of the order of 10^{-9} for $m_{Z'} > 2$ TeV. However, a further suppression could be expected due to a potential suppression of the values of the flavor-changing coupling constants. Under the same assumptions, in 331 models without exotic quarks, the $t \to c\gamma$ branching ratio would receive the dominant contribution from the neutral scalar boson, which could be of the order of 10^{-7} for a Higgs mass of a few hundreds of GeVs.

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I. INTRODUCTION

Despite recent evidence of a neutral Higgs-like particle at the CERN LHC [1], it is necessary to search for effects beyond the Standard Model (SM) as there are still some open questions. Along these lines, several SM extensions have been proposed, such as two-Higgs doublet models [2], left-right symmetric models [3], supersymmetric models [4], left-right supersymmetric models [5], 331 models [6,7], little Higgs models [8,9], and extra dimension models [10], just to mention some of the most popular ones. Such models predict new physics effects in the form of new particles, corrections to the SM couplings, or non-SM couplings. Among the new predicted particles there are, for instance, exotic quarks, CP-even and CP-odd neutral scalar bosons, singly and doubly charged scalar bosons, extra neutral gauge bosons, singly and doubly charged gauge bosons, etc. It may be that there was not enough energy to directly produce any of these hypothetical states at particle colliders, and so their only observable sign would arise indirectly via their loop effects. In particular, the new particles may give rise to sizeable effects on one-loop induced processes, such as the flavor-changing neutral current (FCNC) decays of the top quark $t \to cV$ ($V = \gamma, Z$). Because of its heavy mass, it has been long conjectured that top quark physics offers an opportunity to test the SM and search for new physics effects [11]. The rate for the decay $t \to c\gamma$ is negligibly small in the SM due to the Glashow-Iliopoulos-Maiani mechanism: the respective branching fraction is of the order of 10^{-10} [12]. Since the sensitivity of ATLAS to the $t \to c\gamma$ branching ratio (BR) at the LHC is expected to be of the order of 10^{-4}, it is worth studying such a process in SM extensions, where its branching ratio can be enhanced by several orders of magnitude [13]: this decay has been studied, for instance, in the two-Higgs doublet model [$\mathrm{BR}(t \to c\gamma) \sim 10^{-7}$] [14], technicolor [$\mathrm{BR}(t \to c\gamma) \sim 10^{-7}$] [15], topcolor assisted technicolor [$\mathrm{BR}(t \to c\gamma) \sim 10^{-7}$] [16], supersymmetric models [$\mathrm{BR}(t \to c\gamma) \sim 10^{-6} - 10^{-5}$] [17,18], left-right supersymmetrical models [$\mathrm{BR}(t \to c\gamma) \sim 10^{-6}$] [19], extra dimensions [$\mathrm{BR}(t \to c\gamma) \sim 10^{-10}$] [20], models with an extra neutral gauge boson [$\mathrm{BR}(t \to c\gamma) \sim 10^{-8}$] [21], etc. On the other hand, a model-independent analysis via the effective Lagrangian approach [22] put the upper constraint $\mathrm{BR}(t \to c\gamma) \lesssim 10^{-2}$ by using the experimental bounds on the $b \to s\gamma$ decay. Also, by means of the effective Lagrangian approach, the contribution of a neutral scalar to $\mathrm{BR}(t \to c\gamma)$ was found to be the order of 10^{-8} [23].

We will calculate the $t \to c\gamma$ decay in the framework of models based on the $SU_c(3) \times SU_L(3) \times U_X(1)$ gauge symmetry, which for short are called 331 models and

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have been the source of considerable attention in the literature. The idea of embedding the $SU_L(2) \times U_Y(1)$ gauge group into $SU_L(3) \times U_X(1)$ in order to explain the observation of neutrino-induced trimuon events [24] was discussed in [25], though similar models had already been conjectured [26]. Although these models were soon ruled out, another SM extension based on the $SU_L(3) \times U_X(1)$ gauge group was proposed by the authors of Ref. [6], motivated by the need of a doubly charged gauge boson to restore the unitarity of the cross section of the process $e^- e^- \to W^- V^-$. An almost identical model was proposed independently in [7], but with a different motivation: the need of a chiral theory for doubly charged gauge bosons. Such exotic particles had first been predicted in an $SU(15)$ grand unified theory, which ensured proton stability but required mirror fermions to cancel anomalies [27]. In 331 models, one fermion family must transform under the $SU_L(3)$ group differently from the other two families in order to cancel anomalies, thereby allowing for a solution to the flavor problem: it is necessary that the number of fermion families is a multiple of the quark color number. Also, if the third fermion family is the chosen one to transform differently, 331 models may provide a hint for an eventual understanding of the heaviness of the top quark. Another appealing aspect of these models is that they can accommodate naturally the Peccei-Quinn symmetry [28]. In the phenomenological side, since the $SU_L(2)$ fermion doublets are promoted to $SU_L(3)$ triplets, 331 models require new fermion particles. The specific $SU_L(3)$ representations of the fermion triplets give rise to distinct 331 models. In particular, the most popular ones are the minimal 331 model [6,7] and the 331 model with right-handed neutrinos [29,30]. Other proposed 331 models can be found in Refs. [31–35], and a general treatment of 331 models without exotic quarks can be found in Refs. [36,37]. In addition, although with different purpose and structure, a little Higgs model with global symmetry under the group $[SU(3) \times U(1)]^2$ and local symmetry under the subgroup $SU_L(3) \times U_X(1)$ was proposed in Ref. [38], while its ultraviolet completion was studied in Ref. [39]. This model is an effective theory valid up to the scale of the TeVs, which is known as the simplest little Higgs model and shares the same mechanism of anomaly cancellation as that of the 331 model with right-handed neutrinos.

Apart from reproducing the SM, 331 models predict several new particles. In the gauge sector, the typical signatures are an extra neutral gauge boson and a new singly charged gauge boson. Depending on the particular version of the model, there could be either a new doubly charged gauge boson, as in the minimal 331 model, or a new neutral not self-conjugate gauge boson, as in the 331 model with right-handed neutrinos. As far as the scalar sector is concerned, although the minimal 331 model requires three scalar triplets to accomplish the spontaneous symmetry breaking (SSB) and a sextet to endow the leptons with realistic masses, other versions require a more economical set of scalar multiplets. In this sector there could be new neutral, singly, and doubly charged physical scalar bosons. In the quark sector, three new quarks must be introduced to complete the $SU_L(3)$ triplets: in the minimal 331 model there are three new exotic quarks, two of them have electric charge of $-4/3e$, while the remaining one has charge of $5/3e$; however, there are 331 models in which the new quarks do not have exotic charges [29,31,32,36,37]. The fact that the fermion families transform differently under the gauge group gives rise to FCNCs at the three level mediated by the extra neutral gauge boson and the new neutral scalar bosons, which in turn can induce at the one-loop level the $t \to c\gamma$ decay, which can also be induced by the charged gauge and scalar bosons. Below we will calculate such a decay in the framework of 331 models and analyze the magnitude of the corresponding branching ratio considering the current constraints on the model parameters from experimental data.

The rest of the article is organized as follows. In Sec. II we present an overview of 331 models and their potential sources of flavor change. Section III is devoted to the calculation of the $t \to c\gamma$ decay amplitude, while the numerical analysis and discussion are presented in Sec. IV. The conclusions and outlook are presented in Sec. V.

II. THE MODEL

The motivation and general description of 331 models have already been discussed. We turn to discuss briefly those aspects relevant for our calculation. In 331 models, the charge operator is defined by $Q = T^3 + \beta T^8 + X$, where $T^i = \lambda^i/2$, with λ^i the Gell-Mann matrices and X the $U_X(1)$ quantum number. Specific values of β give rise to distinct models with a peculiar particle content: the minimal 331 model arises when $\beta = \pm\sqrt{3}$, in which case there are three new exotic quarks and a new doubly charged gauge boson; when $\beta = \pm 1/\sqrt{3}$, there are no exotic quarks but new SM-like quarks and a not self-conjugate neutral gauge boson.

The generic contributions to the $t \to c\gamma$ decay in 331 models arise from the extra neutral gauge boson and the neutral Higgs bosons. In this work we will focus mainly on the minimal 331 model, which is the most popular version of these models and the one that predicts additional contributions to the $t \to c\gamma$ decay: those mediated by the exotic quarks along with the new charged gauge and scalar bosons. Nevertheless, our results will be rather general and useful to estimate the size of the $t \to c\gamma$ branching ratio in other 331 models.

In the following, we will not discuss about the lepton sector as it is not relevant for the present work. In the quark sector, three new quarks are required to complete the $SU_L(3)$ triplets. In order to cancel anomalies, the first

two quark families transform under $SU_L(3) \times U_X(1)$ as follows:

$$Q_{1,2} = \begin{pmatrix} u_{1,2} \\ d_{1,2} \\ D_{1,2} \end{pmatrix}: (3, -1/3), \qquad \begin{matrix} u^c_{1,2}: (1, -2/3), \\ d^c_{1,2}: (1, +1/3), \\ D^c_{1,2}: (1, +4/3), \end{matrix} \tag{1}$$

with $D_1 = D$ and $D_2 = S$. The numbers inside the parentheses are the $SU_L(3) \times U_X(1)$ quantum numbers. On the other hand, the third quark family transforms as a triplet:

$$Q_3 = \begin{pmatrix} b \\ -t \\ T \end{pmatrix}: (3^*, 2/3), \qquad \begin{matrix} b^c: (1, +1/3), \\ t^c: (1, -2/3), \\ T^c: (1, -5/3). \end{matrix} \tag{2}$$

As already mentioned, as a consequence of this representation, the new exotic quarks have an electrical charge of $Q_{D,S} = -4/3e$ and $Q_T = 5/3e$.

The scalar sector of 331 models has been studied extensively [40–44]. In the minimal model, one triplet ϕ_Y is necessary to break $SU_L(3) \times U_X(1)$ into $SU_L(2) \times U_Y(1)$, and two triplets, $\phi_{1,2}$, are required to break $SU_L(2) \times U_Y(1)$ into $U_{\rm em}(1)$. In addition, one scalar sextet H is required to give realistic masses to the leptons. More recently, it has been noted that SSB can be achieved with only two scalar triplets [35], but the masses of the charged leptons must be generated via nonrenormalizable effective operators. On the contrary, in 331 models without exotic charge quarks, a scalar sector with two or three scalar triplets is enough to achieve SSB and endow all the particles with masses [41,43,44].

In the minimal 331 model, the scalar triplets have the following quantum numbers:

$$\phi_Y = \begin{pmatrix} \Phi_Y \\ \phi^0 \end{pmatrix}: (3, 1), \qquad \phi_1 = \begin{pmatrix} \Phi_1 \\ \Delta^- \end{pmatrix}: (3, 0), \qquad \phi_2 = \begin{pmatrix} \tilde{\Phi}_2 \\ \rho^{--} \end{pmatrix}: (3, -1), \tag{3}$$

where $\Phi_i = (\phi_i^+, \phi_i^0)^T$ and $\Phi_Y = (G_Y^{++}, G_Y^+)^T$ are $SU_L(2)$ doublets with hypercharge 1 and 3, respectively, and $\tilde{\Phi}_i = i\tau^2\Phi_i^*$. Here G_Y^{++} and G_Y^+ are the would-be Goldstone bosons associated with new doubly and singly charged gauge bosons, whereas the real and imaginary parts of ϕ^0 correspond to one physical Higgs boson and the would-be Goldstone boson associated with an extra neutral gauge boson, respectively. Also, Δ^- and ρ^{--} are singlets of $SU(2)_L$ with hypercharges -2 and -4, respectively. As for the scalar sextet, it has no significance for this work as it is only necessary to give realistic masses to the leptons and so it does not couple to the quarks.

The covariant derivative in the fundamental representation of $SU_L(3) \times U_X(1)$ can be written as

$$D_\mu = \partial_\mu + ig\frac{\lambda^a}{2}W_\mu^a + ig_X X\frac{\lambda^9}{2}V_\mu, \tag{4}$$

where a runs from 1 to 8, W_μ and V_μ are the $SU_L(3)$ and $U_X(1)$ gauge fields, and $\lambda^9 = \sqrt{2/3}\,\mathrm{diag}(1, 1, 1)$. By matching the gauge coupling constants, it is found that $g_X = \sqrt{6}s_W/\sqrt{1 - 4s_W^2}$, with the usual short-hand notation $s_W = \sin\theta_W$.

In the first stage of SSB, the vacuum expectation value (VEV) of the ϕ_Y triplet, $\phi_{Y0}^\dagger = (0, 0, u/\sqrt{2})$, triggers the breaking of the $SU_L(3) \times U_X(1)$ gauge group into $SU_L(2) \times U_Y(1)$, thereby giving rise to two mass-degenerate singly and doubly charged bosons, which are called bileptons as they carry two units of lepton number. They are given as follows:

$$Y_\mu^- = \frac{1}{\sqrt{2}}(W_\mu^6 + iW_\mu^7), \tag{5}$$

$$Y_\mu^{--} = \frac{1}{\sqrt{2}}(W_\mu^4 + iW_\mu^5). \tag{6}$$

There are also a massive extra neutral gauge boson Z'_μ and a massless gauge boson B_μ, which are given in terms of W_μ^8 and V_μ. While B_μ corresponds to the $U_Y(1)$ gauge field, the massless fields associated with the unbroken generators of $SU_L(3)$, W_μ^i $(i = 1, 2, 3)$, turn out to be the gauge fields of the $SU_L(2)$ group.

At the Fermi scale, the SM gauge group is spontaneously broken down into the electromagnetic group via the VEVs of the $SU_L(2)$ doublets, $\langle\Phi_i^0\rangle_0 = (0, v_i/\sqrt{2})^T$ $(i = 1, 2)$. The SM charged gauge bosons, $W_\mu^\pm = (W_\mu^1 \mp iW_\mu^2)/\sqrt{2}$, get their masses and the bileptons receive additional mass contributions. Finally, the W^3, W^8, and V gauge fields define three neutral fields as follows:

$$\begin{pmatrix} W_\mu^3 \\ W_\mu^8 \\ V_\mu \end{pmatrix} = \begin{pmatrix} s_W & c_W & 0 \\ \sqrt{3}s_W & -\sqrt{3}s_W t_W & -\frac{1}{c_W}\sqrt{1 - 4s_W^2} \\ \sqrt{1 - 4s_W^2} & -\sqrt{1 - 4s_W^2}\,t_W & \sqrt{3}t_W \end{pmatrix} \times \begin{pmatrix} A_\mu \\ Z_\mu \\ Z'_\mu \end{pmatrix}, \tag{7}$$

where A_μ corresponds to the photon, but Z_μ and Z'_μ need to be rotated to obtain the mass eigenstates: the SM neutral weak gauge boson $Z_1 = Z\cos\theta - Z'\sin\theta$ and the extra neutral gauge boson $Z_2 = Z\sin\theta + Z'\cos\theta$, with the mixing angle θ defined by $\sin^2\theta = (m_Z^2 - m_{Z_1}^2)/(m_{Z_2}^2 - m_{Z_1}^2)$. Since θ is strongly constrained by experimental data, we will assume that $\theta \simeq 0$ and thus the Z and Z' gauge bosons will be taken as the mass eigenstates. The masses of the heavy physical states are thus

$$m_{Y^{--}}^2 = \frac{g^2}{4}(u^2 + v_2^2 + 4v_3^2), \tag{8}$$

$$m_{Y^-}^2 = \frac{g^2}{4}(u^2 + v_1^2 + v_3^2), \tag{9}$$

$$m_{Z'}^2 = \frac{g^2}{3(1-4s_W^2)} \times \left(c_W^2 u^2 + \frac{(1-4s_W^2)^2}{4c_W^2}(v_1^2 + v_2^2 + v_3^2) + 3s_W^2 v_2^2\right), \tag{10}$$

where v_3 is the VEV of the Φ_3 doublet, which is required to endow the leptons with masses. From the symmetry breaking hierarchy, $u > v_{1,2} > v_3$, it turns out that $m_{Z'} > m_{Y^-,Y^{--}} > m_{W,Z}$. In fact, neglecting the splitting between the bilepton masses, $m_{Y^{--}} \simeq m_{Y^-} \equiv m_Y$, we obtain the following approximate relation [45]:

$$m_Y \simeq \sqrt{\frac{3}{4}}\frac{\sqrt{1-4s_W^2}}{c_W} m_{Z'} \simeq \frac{1}{3} m_{Z'}. \tag{11}$$

After SSB and once the gauge eigenstates are rotated to mass eigenstates, the Yang-Mills Lagrangian for the fields W^μ and V^μ can be decomposed into the SM Yang-Mills Lagrangian plus a term that contains the interactions between the SM gauge bosons and the heavy charged gauge bosons, together with a term that only contains interactions between the Z' gauge boson and the bileptons. The term necessary for our calculation can be written as

$$\mathcal{L}^{\mathrm{SM-331}} = -\frac{1}{2}(D_\mu Y_\nu - D_\nu Y_\mu)^\dagger (D^\mu Y^\nu - D^\nu Y^\mu) - Y^{\dagger\mu}(ig\mathbf{W}_{\mu\nu} + ig'\mathbf{B}_{\mu\nu})Y^\nu, \tag{12}$$

where $\mathbf{W}_{\mu\nu} = \tau^i W^i_{\mu\nu}/2$, $\mathbf{B}_{\mu\nu} = YB_{\mu\nu}/2$, and $Y_\mu = (Y_\mu^{--}, Y_\mu^-)^T$; also, $D_\mu = \partial_\mu - ig\mathbf{W}_\mu - ig'\mathbf{B}_\mu$ is the covariant derivative associated with the electroweak group. From here we can get the interactions of the bilepton gauge bosons with the photon.

As to the neutral and charged currents mediated by the heavy gauge bosons, they arise from the fermion kinetic terms and can be written as

$$\mathcal{L}^{\bar{q}'q'V} = -\frac{g}{2c_W}\left(\sum_{i=1}^{3} \bar{Q}'_{Li}\gamma^\mu \mathbf{H}_\mu Q'_{Li} + \sum_{i=1}^{9} 6s_W^2 Z'_\mu \bar{q}'_{Ri}\gamma^\mu X q'_{Ri}\right), \tag{13}$$

with

$$\mathbf{H}_\mu = \begin{pmatrix} \left(\frac{2(3X+2)s_W^2-1}{\sqrt{3}\sqrt{1-4s_W^2}}\right)Z'_\mu & 0 & \sqrt{2}c_W Y_\mu^{--} \\ 0 & \left(\frac{2(3X+2)s_W^2-1}{\sqrt{3}\sqrt{1-4s_W^2}}\right)Z'_\mu & \sqrt{2}c_W Y_\mu^{-} \\ \sqrt{2}c_W Y_\mu^{++} & \sqrt{2}c_W Y_\mu^{+} & 2\left(\frac{(3X-4)s_W^2+1}{\sqrt{3}\sqrt{1-4s_W^2}}\right)Z'_\mu \end{pmatrix}, \tag{14}$$

where Q'_i ($i = 1, 2, 3$) is a quark triplet and q'_i is a quark singlet, both in the flavor basis. Since the third family has a different representation under $SU_L(3)$, after the flavor eigenstates are rotated to the mass eigenstates there emerge FCNC couplings mediated by the Z' gauge boson. The flavor conserving Z' couplings to a quark pair have the form

$$\mathcal{L}^{\bar{q}qZ'} = -\frac{g}{c_W} Z'_\mu \bar{q}\gamma^\mu (g_L'^q P_L + g_R'^q P_R) q, \tag{15}$$

where $P_{L,R}$ are the chiral projection operators and the $g_{L,R}'^q$ constants are presented in the Appendix. On the other hand, the flavor-changing neutral and charged currents required by our calculation can be arranged as [46]

$$\mathcal{L}^{NC} = -\frac{gc_W}{\sqrt{3}\sqrt{1-4s_W^2}} Z'_\mu U^*_{L3i} U_{L3j} \bar{u}_i \gamma^\mu P_L u_j, \tag{16}$$

$$\mathcal{L}^{CC} = -\frac{g}{\sqrt{2}}(Y_\mu^- U^*_{Li3}\bar{u}_i\gamma^\mu P_L T + Y_\mu^{--}(U_{Li1}\bar{D}\gamma^\mu P_L u_i + U_{Li2}\bar{S}\gamma^\mu P_L u_i)) + \mathrm{H.c.}, \tag{17}$$

where U_L, which is the 3×3 matrix that diagonalizes the SM up quarks from flavor eigenstates to mass eigenstates, is related to the Cabibbo-Kobayashi-Maskawa (CKM) matrix by $U_{\mathrm{CKM}} = U_L^\dagger V_L$, with V_L the mass matrix that diagonalizes the SM down quarks. Notice that the $D_{1,2}$ flavor eigenstates can be chosen as the mass eigenstates since the two first fermion families transform symmetrically [46].

Finally, we will discuss briefly about the Yukawa couplings associated with the quark sector, which can be written in terms of the SM quark doublets $q'_i = (u'_i, d'_i)^T$ as [46]

$$-\mathcal{L} = \sum_{k=1}^{3}\sum_{i=1}^{2}(\bar{q}'_{Li}h_d^{ik}d'_{Rk}\Phi_1 + \bar{q}'_{Li}h_u^{ik}u'_{Rk}\tilde{\Phi}_2) + \sum_{k=1}^{3}(\bar{q}'_{L3}h_d^{3k}d'_{Rk}\Phi_2 + \bar{q}'_{L3}h_u^{3k}u'_{Rk}\tilde{\Phi}_1)$$
$$+ \sum_{k=1}^{3}\sum_{i=1}^{2}(\bar{D}'_{Li}h_d^{ik}d'_{Rk}\Delta^- + \bar{D}'_{Li}h_u^{ik}u'_{Rk}\rho^{--}) + \sum_{k=1}^{3}(\bar{T}_L h_d^{3k}d'_{Rk}\rho^{++} - \bar{T}_L h_u^{3k}u'_{Rk}\Delta^+)$$
$$+ \sum_{i=1}^{2}\sum_{j=1}^{2}(\bar{D}'_{Li}h_D^{ij}D'_{Rj}\phi^0 + \bar{q}'_{Li}h_D^{ij}D'_{Rj}\phi_Y) + \bar{T}_L h_T T_R \phi^{0*} - \bar{q}'_{L3}h_T T_R \tilde{\phi}_Y + \text{H.c.}, \quad (18)$$

where h^{ij} are symmetric matrices in flavor space. After the first stage of SSB, there are two-Higgs doublets plus one neutral, one singly charged, and one doubly charged scalar bosons. There will be additional scalar multiplets, which arise from the scalar sextet, that do not couple to the quarks. We can observe from the last line of Eq. (18) that after the ϕ_Y doublet develops a VEV, the exotic quarks get their masses, which are thus of the order of u. Furthermore, after SSB and once the mass eigenstates are obtained, there is a plethora of physical Higgs bosons (5 neutral CP-even, 3 neutral CP-odd, 4 singly charged, and 3 doubly charged scalar bosons), which are a mix of the Higgs eigenstates [40,42,47]. Since these physical Higgs bosons can induce flavor change, they will contribute to the decay $t \to c\gamma$. In particular there could be flavor change mediated by neutral scalar Higgs bosons due to the asymmetry in the $SU(3)$ representation of the quark families. However, the introduction of an *ad hoc* discrete symmetry can eliminate any dangerous FCNC. Since a complete treatment of the scalar sector is rather complicated and requires us to consider several parameters, we will take instead a more practical approach: we will consider a dimension-four effective Lagrangian for typical neutral and charged scalars that can induce the $t \to c\gamma$ decay. For the effective couplings we can consider the so-called Cheng-Sher ansatz [48], which is suited for models with multiple Higgs doublets. This will be useful to estimate the size of the potential contributions to the $t \to c\gamma$ branching ratio in 331 models.

III. DECAY $t \to c\gamma$ IN 331 MODELS

We find it useful to present our results in a model-independent fashion. We thus consider the following renormalizable interactions that can induce the $t \to c\gamma$ decay. We start with the interactions between a neutral Higgs boson ϕ and a quark pair:

$$\mathcal{L}^{\phi} = -\frac{g}{c_W}\bar{q}_i(L_\phi^{ij}P_L + R_\phi^{ij}P_R)q_j\phi, \quad (19)$$

where $i, j = 1, 2, 3$ stands for the quark flavors, while L_ϕ^{ij} and R_ϕ^{ij} are coupling constants. If CP is conserved then $L_\phi^{ij} = R_\phi^{*ji}$. From now on, unless stated otherwise, ϕ will denote a neutral scalar boson. As for the singly and doubly charged scalars, ϕ^- and ϕ^{--}, their interactions with SM up quarks and the exotic quarks, D_i and T, can be expressed in the form

$$\mathcal{L}^{\text{SCC}} = -\frac{g}{c_W}\bar{u}_i(L_{\phi^-}^{iT}P_L + R_{\phi^-}^{iT}P_R)T\phi^-$$
$$-\frac{g}{c_W}\sum_{i=1,2}\bar{D}_i(L_{\phi^{--}}^{ij}P_L + R_{\phi^{--}}^{ij}P_R)u_j\phi^{--} + \text{H.c.} \quad (20)$$

As far as the gauge sector is concerned, the most general renormalizable interactions of a neutral gauge boson Z' with a quark pair can be written as:

$$\mathcal{L}^{Z'} = -\frac{g}{c_W}\bar{q}_i\gamma^\mu(L_{Z'}^{ij}P_L + R_{Z'}^{ij}P_R)q_jZ'_\mu. \quad (21)$$

Finally, the interactions of the singly and doubly charged gauge bosons to SM and exotic quark are

$$\mathcal{L}^{\text{GCC}} = -\frac{g}{c_W}\bar{u}_i\gamma^\mu(L_{Y^-}^{iT}P_L + R_{Y^-}^{iT}P_R)TY_\mu^-$$
$$-\frac{g}{c_W}\sum_{i=1,2}\bar{D}_i\gamma^\mu(L_{Y^{--}}^{ij}P_L + R_{Y^{--}}^{ij}P_R)u_jY_\mu^{--} + \text{H.c.} \quad (22)$$

We also need the interactions with the photon, which are dictated by electrodynamics and follow from Eq. (12) and the kinetic term of the scalar mutiplets. These interactions and all the Feynman rules necessary for our calculation are presented in the Appendix.

Because of electromagnetic gauge invariance, the $t(p_1) \to c(p_2)\gamma(q)$ decay amplitude can be cast in the form

$$\mathcal{M}(t \to c\gamma) = \frac{i}{m_t}\bar{c}(p_2)\sigma_{\mu\nu}(C_LP_L + C_RP_R)t(p_1)q^\nu\epsilon^\mu(q). \quad (23)$$

We show in Fig. 1 the one-loop contributions to the $t \to c\gamma$ decay from an arbitrarily charged gauge boson V and a SM or exotic quark. We are using the unitary gauge, so there are no contributions from nonphysical particles. There are also bubble diagrams that can contribute to the on-shell $\bar{c}t\gamma$ vertex. Although such diagrams do not contribute to the dipole coefficients $C_{L,R}$, they give rise to ultraviolet divergent terms that violate electromagnetic gauge invariance, which we have verified are canceled out by similar terms arising from the triangle diagrams. This is similar to what happens with the $b \to s\gamma$ decay [49]. There are four possible combinations of loops carrying a quark and a gauge boson with the following electric charges: (i) $Q_q = 5/3e$ and $Q_V = e$,

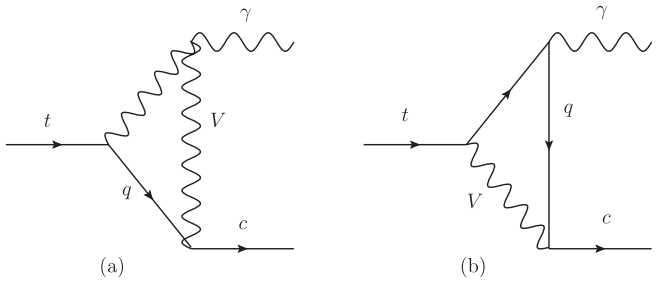

FIG. 1. Feynman diagrams, in the unitary gauge, for the decay $t \to c\gamma$ in 331 models. We show the contribution of an arbitrarily charged gauge boson V along with an exotic or SM-like quark. The diagram (a) does not contribute in the case of the neutral gauge boson. The contribution of charged scalar particles are similar but with the gauge boson replaced by the scalar boson. There are also bubble diagrams that do not contribute to the $t \to c\gamma$ amplitude but are necessary to render electromagnetic gauge invariance and cancel out ultraviolet divergences.

(ii) $Q_q = -4/3e$ and $Q_V = -2e$, (iii) $Q_q = 2/3e$ and $Q_V = 0$, and (iv) $Q_q = -1/3e$ and $Q_V = -e$. As for the contribution of an arbitrarily charged scalar boson, it arises from similar diagrams to that shown in Fig. 1 but with the gauge boson replaced by a scalar boson.

We now turn to present the results for the contributions to the $t \to c\gamma$ decay arising from the interactions (19)–(22). In order to solve the one-loop tensor integrals, we expressed them in terms of scalar two- and three-point scalar functions via the Passarino-Veltman reduction scheme [50]. Analytical results are given for these scalar functions in terms of dilogarithms and other transcendental functions.

A. Gauge boson contribution

We denote the coupling constants appearing in Eqs. (21) and (22) by L_V^{ij} and R_V^{ij}. After a lengthy algebra, we obtain the contribution from the loops carrying the arbitrarily charged gauge boson V and the quark q:

$$
\begin{aligned}
C_R^V = \frac{3eg^2}{16\pi^2 c_W^2}\frac{1}{2\delta_{tc}}\Big(&\tilde{\Delta}_{qt}\sqrt{x_t}((\sqrt{x_c}(\delta_{qt}+2)R_V^{tq}+\sqrt{x_q}\delta_{tc}L_V^{tq})R_V^{qc}-\sqrt{x_t}(\delta_{qc}+2)L_V^{qc}L_V^{tq})\\
&-\frac{Q_t\delta_q}{\sqrt{x_c}}(\sqrt{x_c}(\delta_{qt}+2)L_V^{qc}L_V^{tq}+\sqrt{x_q}\delta_{tc}L_V^{qc}R_V^{tq}-\sqrt{x_t}(\delta_{qc}+2)R_V^{tq}R_V^{qc})G_{(c,q,V)}+\frac{2Q_q}{\sqrt{x_q}}(\sqrt{x_c x_t x_q}(\delta_{qc}+2)R_V^{tq}R_V^{qc}\\
&-x_t(\sqrt{x_c}\delta_{tc}R_V^{tq}+\sqrt{x_q}(\delta_{qt}+2)L_V^{tq})L_V^{qc})H_{(c,t,q,V)}+2\Delta_{qt}(\sqrt{x_t}(\sqrt{x_c}(\delta_{qc}+\delta_{tc}+2)R_V^{tq}-3\sqrt{x_q}\delta_{tc}L_V^{tq})R_V^{qc}\\
&-x_t(\delta_{qt}-\delta_{tc}+2)L_V^{qc}L_V^{tq})H_{(c,t,V,q)}+\frac{1}{\delta_{tc}}[x_t\sqrt{x_c}(Q_t(x_c(x_q-3)+\delta_q(x_q+2))-x_t(\delta_{qc}\tilde{\Delta}_{qt}+4\Delta_{qt}))R_V^{qc}R_V^{tq}\\
&+(Q_t(2x_t(x_q^2-(1-\delta_q)\delta_t)-x_c(x_q(\delta_{qt}-\delta_t)-\delta_t-1))+2\Delta_{qt}x_c(x_q+2)x_t-\tilde{\Delta}_{qc}x_c x_t^2)L_V^{qc}L_V^{tq}\\
&-\sqrt{x_c x_q}\delta_{tc}(x_t(\Delta_{qt}-3Q_q)-Q_t\delta_q)L_V^{qc}R_V^{tq}-6\sqrt{x_t x_q}Q_t\delta_{tc}R_V^{qc}L_V^{tq}]F_{(c,t,q,V)}\Big),
\end{aligned}
\tag{24}
$$

and

$$
C_L^V = C_R^V(t \leftrightarrow c), \tag{25}
$$

where $\Delta_{ij} = Q_i - Q_j$ and $\tilde{\Delta}_{ij} = 2Q_i - Q_j$, with Q_i the electric charge of particle i in units of e; we also introduced the definitions $x_i = m_i^2/m_V^2$, $\delta_{ij} = x_j - x_i$, and $\delta_i = 1 - x_i$. The $F_{(i,j,k,l)}$, $G_{(i,j,k)}$, and $H_{(i,j,k,l)}$ functions are given in terms of Passarino-Veltman scalar functions:

$$
F_{(i,j,k,l)} = B_0(m_i^2, m_k^2, m_l^2) - B_0(m_j^2, m_k^2, m_l^2), \tag{26}
$$

$$
G_{(i,j,k)} = B_0(0, m_j^2, m_k^2) - B_0(m_i^2, m_j^2, m_k^2), \tag{27}
$$

$$
H_{(i,j,k,l)} = m_k^2 C_0(m_i^2, m_j^2, 0, m_k^2, m_l^2, m_k^2), \tag{28}
$$

where B_0 and C_0 are two- and three-point scalar functions written in the notation of Ref. [51]. From here it is clear that ultraviolet divergences cancel out. Explicit integration of the above functions yields

$$
\begin{aligned}
F_{(i,j,k,l)} = f(m_i^2, m_j^2, m_k^2, m_l^2) &+ \frac{1}{m_i^2}(f_- - f_+)(m_i^2, m_k^2, m_l^2)\\
&- \frac{1}{m_j^2}(f_- - f_+)(m_j^2, m_k^2, m_l^2),
\end{aligned}
\tag{29}
$$

$$
G_{(i,j,k)} = g(m_i^2, m_j^2, m_k^2) + \frac{1}{m_i^2}(f_+ - f_-)(m_i^2, m_j^2, m_k^2), \tag{30}
$$

$$
\begin{aligned}
H_{(i,j,k,l)} = \frac{m_k^2}{m_j^2 - m_i^2}&((h_+ + h_-)(m_i^2, m_k^2, m_l^2)\\
&- (h_+ + h_-)(m_j^2, m_k^2, m_l^2)),
\end{aligned}
\tag{31}
$$

with

$$f(x,y,w,z) = \frac{(x-y)(w-z)}{2xy}\log\left(\frac{w}{z}\right), \quad (32)$$

$$f_\pm(x,y,z) = \sqrt{\lambda(x,y,z)}\,\mathrm{arctanh}\left(\frac{y-z\pm x}{\sqrt{\lambda(x,y,z)}}\right), \quad (33)$$

$$g(x,y,z) = \frac{(y-z)^2 - x(y+z)}{2x(y-z)}\log\left(\frac{y}{z}\right) - 1, \quad (34)$$

$$h_\pm(x,y,z) = \mathrm{Li}_2\left(\frac{2x}{x+y-z\pm\sqrt{\lambda(x,y,z)}}\right), \quad (35)$$

and $\lambda(x,y,z) = (x-y-z)^2 - 4yz$.

A special case arises when only the left-handed quarks interact with the exchanged gauge boson, as occurs with the bilepton contributions in the minimal 331 model. In this scenario we have

$$C_R^Y = \frac{3eg^2}{16\pi^2 c_W^2}\frac{L_Y^{tq}L_Y^{qc}}{2\delta_{tc}}\Big(\tilde{\Delta}_{qt}x_t(\delta_{qc}+2) + \frac{1}{\sqrt{x_c}}Q_t\delta_q(\sqrt{x_c}\delta_{qt}+2)G_{(c,q,Y)} + 2Q_q x_t(\delta_{qt}+2)H_{(c,t,q,Y)}$$
$$- 2\Delta_{qt}x_t(\delta_{qt}-\delta_{tc}+2)H_{(c,t,Y,q)} - \frac{1}{\delta_{tc}}(Q_t(2x_t(x_q^2-(1-\delta_q)\delta_t) - x_c(x_q(\delta_{qt}-\delta_t)-\delta_t-1))$$
$$+ 2\Delta_{qt}x_c(x_q+2)x_t - \tilde{\Delta}_{qc}x_c x_t^2)F_{(c,t,q,Y)}\Big). \quad (36)$$

B. Scalar contribution

We now consider the contribution from the loops carrying an arbitrarily charged scalar boson ϕ and a quark q, which arise from Feynman diagrams similar to those of Fig. 1 but with the gauge boson replace by a scalar boson. The Passarino-Veltman reduction scheme yields the following $C_{L,R}^\phi$ coefficients:

$$C_R^\phi = \frac{3eg^2}{16\pi^2 c_W^2}\frac{1}{2\eta_{tc}}\Big(\frac{\sqrt{y_t}}{\sqrt{y_q}}Q_q(\sqrt{y_q}(\sqrt{y_c}L_\phi^{qc}R_\phi^{tq} - \sqrt{y_t}L_\phi^{tq}R_\phi^{qc}) - \eta_{tc}R_\phi^{tq}R_\phi^{qc})H_{(c,t,q,\phi)} + 2\sqrt{y_t}(\sqrt{y_c}L_\phi^{qc}R_\phi^{tq} - \sqrt{y_t}L_\phi^{tq}R_\phi^{qc})$$
$$\times(\tilde{\Delta}_{qt} + \Delta_{qt}H_{(c,t,\phi,q)}) + \frac{1}{\sqrt{y_c}}Q_t\eta_q(\sqrt{y_c}L_\phi^{tq}R_\phi^{qc} - \sqrt{y_t}L_\phi^{qc}R_\phi^{tq})G_{(c,q,\phi)} + \frac{1}{\eta_{tc}}((y_c y_t\tilde{\Delta}_{qt} + Q_t\eta_q(y_c-2y_t))L_\phi^{tq}R_\phi^{qc}$$
$$+ \sqrt{y_t}(\sqrt{y_c}(Q_t\eta_q - y_t\tilde{\Delta}_{qt})L_\phi^{qc} + 2Q_t\sqrt{y_q}\eta_{tc}R_\phi^{qc})R_\phi^{tq})F_{(c,t,q,\phi)}\Big), \quad (37)$$

where $y_i = m_i^2/m_\phi^2$, $\eta_{ij} = y_i - y_j$, and $\eta_i = 1 - y_i$. C_L^ϕ can be obtained from C_R^ϕ after the replacements $t \leftrightarrow c$, $L_\phi^{tq} \leftrightarrow R_\phi^{qc}$, and $R_\phi^{tq} \leftrightarrow L_\phi^{qc}$ are done:

$$C_L^\phi = C_R^\phi\left(\begin{array}{c} t \leftrightarrow c \\ L \leftrightarrow R \end{array}\right). \quad (38)$$

For a neutral CP-even neutral scalar boson, denoted by ϕ rather than ϕ^0 to avoid being plagued by indices, $R_\phi^{ij} = L_\phi^{ij} \equiv \lambda_\phi^{ij}$, the above expression reduces to

$$C_R^\phi = \frac{eg^2}{16\pi^2 c_W^2}\frac{\sqrt{y_t}\lambda_\phi^{qc}\lambda_\phi^{tq}}{(\sqrt{y_c}+\sqrt{y_t})}\Big(1 + \frac{1}{\sqrt{y_c}}\eta_q G_{(c,q,\phi)} + \frac{2}{\sqrt{y_q}}(\sqrt{y_q}+\sqrt{y_c}+\sqrt{y_t})H_{(c,t,q,\phi)}$$
$$+ \frac{1}{\eta_{tc}}\Big(\sqrt{y_c y_t} - 2\sqrt{y_q}(\sqrt{y_c}+\sqrt{y_t}) + \eta_q\Big(2+\sqrt{\frac{y_c}{y_t}}\Big)\Big)F_{(c,t,q,\phi)}\Big). \quad (39)$$

IV. NUMERICAL ANALYSIS AND DISCUSSION

In addition to the SM contribution to the $t \to c\gamma$ decay, the new contribution from the gauge sector of the minimal 331 model can be written as

$$C_{L,R}^G = \sum_{q=u,c,t} C_{L,R}^{Z'} + \sum_{q=T} C_{L,R}^{Y^-} + \sum_{q=D,S} C_{L,R}^{Y^{--}}, \quad (40)$$

while the contribution from the Higgs sector is as follows:

$$C_{L,R}^H = \sum_{q=u,c,t} C_{L,R}^{\phi} + \sum_{q=T} C_{L,R}^{\phi^-} + \sum_{q=D,S} C_{L,R}^{\phi^{--}}, \quad (41)$$

where it is assumed that we must consider all the physical neutral and charged Higgs bosons. In the case of 331 models without exotic quarks, the generic contribution arises only from the Z' gauge boson and the neutral scalar bosons.

From Eq. (23), the corresponding $t \to c\gamma$ decay width follows easily:

$$\Gamma(t \to c\gamma) = \frac{m_t}{16\pi}\left(1 - \frac{m_c^2}{m_t^2}\right)^3 (|C_L^G + C_L^H|^2 + |C_R^G + C_R^H|^2). \tag{42}$$

In order to get a realistic estimate for $\Gamma(t \to c\gamma)$, we will consider the current constraints on the masses of the heavy gauge boson, the exotic quarks, and the scalar bosons.

A. Constraints on the model parameters

Considerable work has gone into studying constraints on the masses of the extra gauge bosons of 331 models. The most stringent bound on the mass of a doubly charged bilepton was obtained from the data of the search for muonium to antimuonium conversion [52]. This bound, $m_{Y^{--}} > 800$ GeV, is based on the assumptions that the bilepton-lepton couplings are flavor diagonal and the scalar sector of the model does not contribute significantly to muonium-antimuonium conversion. Another stringent bound, $m_{Y^{--}} > 750$ GeV, arises from fermion pair production and lepton-flavor violating processes [53]. It has been argued [54], however, that these bounds can be evaded if one makes less restrictive assumptions than the aforementioned analyses. As for the Z' gauge boson mass, it is related to the bilepton masses by Eq. (11): $m_{Z'} \simeq 3m_Y$. Therefore, the most stringent bounds on the doubly charged bilepton mass translates into a lower bound on $m_{Z'}$ of about 2 TeV, which is similar to other restrictive bounds obtained in Refs. [55–57].

There is considerably less literature dealing with bounds on the exotic quark masses [58,59], which, in general, depend on the masses of the heavy gauge bosons. From the search for supersymmetric particles at the Tevatron, a bound on the D quark mass of about 300 GeV was obtained for $m_{Z'}$ around 1 TeV [58]. Another constraint was obtained in [59] from the experimental data on the $Z \to \bar{b}b$ decay and electroweak precision measurements at the Z pole: it was found that the T quark mass is bounded into the 1500–4000 GeV interval for $m_{Y^{--}}$ around 700 GeV.

As far as the bounds on the scalar boson masses are concerned, these are more difficult to obtain as the scalar sector of the minimal 331 model is plagued with free parameters. There are a few recent bounds on the charged scalar boson masses from direct searches at the LHC, but they are model dependent. As will be discussed below, we will consider the scenario in which there is only a relatively light neutral scalar boson, with a mass of a few hundreds of GeVs, whereas the remaining scalar bosons will be assumed to be very heavy. Hence, the bulk of the scalar contribution would arise from the neutral scalar boson.

In conclusion, in our analysis below we will consider degenerate bileptons with a mass above 600 GeV, whereas for the extra neutral gauge boson mass we will consider the relation $m_{Z'} \simeq 3m_Y$. For the exotic quarks we will assume the hierarchy $m_T \sim m_S > m_D$, with m_D, and m_S around 500 and 1000 GeV, respectively. Furthermore, the existence of a relatively light neutral Higgs boson with FCNC couplings and a mass of a few hundreds of GeVs will be assumed.

B. Gauge boson contribution

In order to estimate the size of the $t \to c\gamma$ decay width, we show in Fig. 2 the behavior of the partial contributions from the heavy gauge bosons to the C_R^V coefficients as a function of the gauge boson and the exotic quark masses. We do not show the C_L^V coefficient as its size is more than two orders of magnitude below than that of C_R^V due to the small value of the c mass. Each contribution was divided by the associated products of U_L matrix elements, which are encapsulated in the η coefficient. A word of caution is in order here: the values shown in Fig. 2 can be dramatically reduced if η is much smaller than unity. In the case of the extra neutral gauge boson Z', we only show the contribution from the loops carrying the c and t quarks since the amplitude corresponding to the u quark involves two flavor-changing vertices and it is expected to be more suppressed. From Fig. 2, we can conclude that the largest contribution to C_R^V could arise from the charged bileptons, while the smallest contribution could be due to the extra neutral gauge boson. Another point worth mentioning is that, while the C_R^V coefficient is strongly dependent on the value of the gauge boson mass and can decrease up to 1 order of magnitude in the interval from 600 to 2000 GeV, its change is almost imperceptible when the exotic quark mass is varied in a similar interval, which is evident in Fig. 2 as the curves corresponding to exotic quarks with same electric charges but distinct masses almost overlap. Although there could be large cancellations when summing over all the contributions to C_R^V, it is interesting to point out that a mechanism such as the one that suppresses

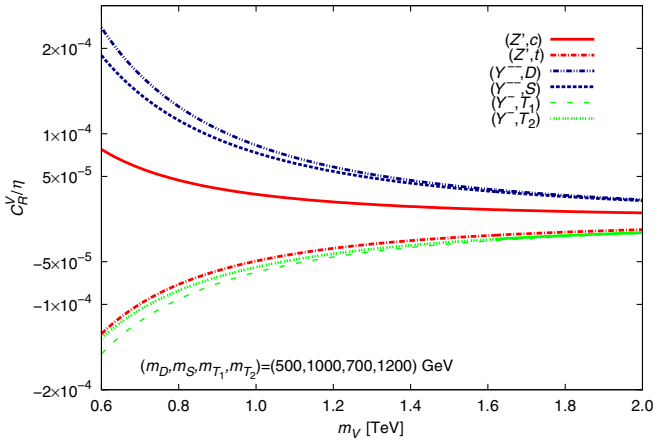

FIG. 2 (color online). Partial contribution from the heavy gauge boson and the internal quark pair (V, q) to C_R^V as a function of m_V and for fixed values of the exotic quark masses in the minimal 331 model. We considered two values for the mass of the T quark (m_{T_1} and m_{T_2}). Also $\eta = U_{L21}U_{L31}^*$ for (Y^{--}, D), $\eta = U_{L22}U_{L32}^*$ for (Y^{--}, S), and $\eta = U_{L23}U_{L33}^*$ for (Z', c), (Z', t), and (Y^-, T).

the contribution from the W gauge boson does not operate in the case of the charged bileptons, even if we assume that they are mass degenerate. In such a case, from Eq. (36) we can see that C_R^Y would adopt the form

$$C_R^Y = \sum_{i=1}^{3} U_{L2i} U_{L3i}^* f(m_{Q_i}, m_Y, Q_Y), \tag{43}$$

where the sum runs over the exotic quarks. Here m_Y stands for the bilepton mass, while the bilepton electric charge is $Q_Y = Q_i - Q_t$. Since U_L is unitary, C_R^Y would vanish if the f function was independent of m_{Q_i} and Q_i. However, even if the exotic quarks were mass degenerate, C_R^Y would not vanish as they do not have the same electric charge. Therefore, we do not expect a strong suppression of C_R^Y. In the case of the Z' contribution, we also do not expect large cancellations between the c and t contributions because of the disparity of the top quark mass and the nonuniversality of the couplings of the Z' to SM quarks.

We now consider a scenario in which the heavy gauge boson contributions to the $t \to c\gamma$ decay add up instead of canceling out. In such a case, a rough estimate for the branching ratio, $\mathrm{BR}(t \to c\gamma)$, is that it would be of the same order of magnitude than its partial contributions: although we would need to know the actual values of the U_L matrix elements to obtain the total contribution, an enhancement of several orders of magnitude with respect to each contribution cannot be expected. We thus show the individual behavior of the partial contributions to $\mathrm{BR}(t \to c\gamma)$ in Fig. 3 as a function of the bilepton mass and illustrative values of the

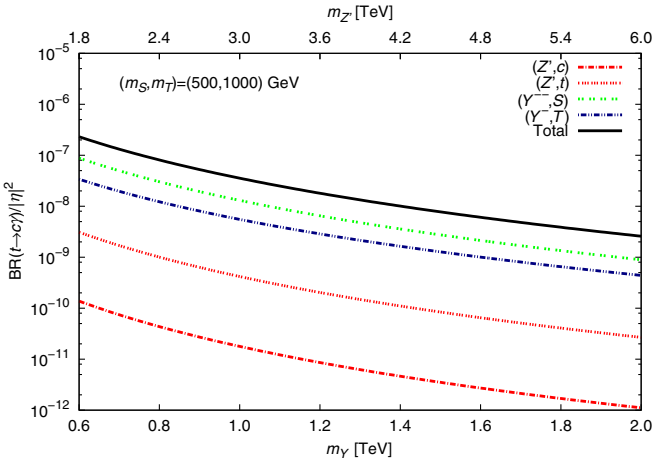

FIG. 3 (color online). Partial contributions from the heavy gauge boson and the internal quark (V, q) to the $t \to c\gamma$ branching ratio in the minimal 331 model as a function of the bilepton mass, considering degenerate bileptons, and for fixed values of the exotic quark masses. The Z' mass is related to m_Y by Eq. (11): $m_{Z'} \simeq 3m_Y$. Also $\eta = U_{L22}U_{L32}^*$ for (Y^{--}, S), and $\eta = U_{L23}U_{L33}^*$ for (Z', c), (Z', t), and (Y^-, T). We also show the total contribution to $\mathrm{BR}(t \to c\gamma)$ considering the scenario discussed in the text, namely, when $U_{L21}U_{L31}^* \sim 0$ and $U_{L22}U_{L32}^* \sim -U_{L23}U_{L33}^*$.

exotic quark masses. Since we use the relation $m_{Z'} \simeq 3m_Y$, the Z' contribution is considerably more suppressed as compared to the bilepton contribution. We can conclude that it would be very unlikely that the total contribution of the heavy gauge bosons to $\mathrm{BR}(t \to c\gamma)$ would surpass the 10^{-7} level even if there were no large cancellations between the partial contributions or a further suppression coming from the U_L matrix elements. In order to illustrate our point, we assume a simple scenario in which $U_{L21}U_{L31}^* \sim 0$ and $U_{L22}U_{L32}^* \sim -U_{L23}U_{L33}^*$. We then calculate the total contribution to the $t \to c\gamma$ branching ratio from the heavy gauge bosons: the result is given by the solid line shown in the plot of Fig. 3. Although the total $\mathrm{BR}(t \to c\gamma)$ is slightly enhanced, even if we consider nondegenerate bileptons, such an enhancement would hardly surpass 1 order of magnitude. Finally, we also expect that the generic contribution to $\mathrm{BR}(t \to c\gamma)$ from the gauge sector of 331 models is of the order of 10^{-9} at most since it only arises from the extra neutral Z' gauge boson whose mass is constrained, from experimental data, to be much larger than 1 TeV.

C. Scalar boson contribution

Since there are several physical neutral, singly, and doubly charged scalars, along with several free parameters, such as the masses of the scalar bosons, mixing angles, and Yukawa couplings, the analysis of this contribution turns out to be very complicated. Fortunately, at low energies, as far as the quark sector is concerned, the scalar sector of the minimal 331 model resembles that of a two-Higgs doublet model [60]. Therefore, in our analysis we will assume that the largest contribution from the scalar sector arises from the lightest neutral scalar boson. This is equivalent to assuming that the remaining scalar bosons are very heavy or that there is a large suppression of the associated Yukawa couplings. Hence we will analyze the behavior of the contribution from a typical neutral scalar boson with a mass of a few hundreds of GeVs. For the coupling of such a scalar boson to a SM quark pair we will consider the Cheng-Sher ansatz [48], which is meant for multiple-Higgs-doublets models. We will thus assume that $\lambda_\phi^{ij} = \sqrt{m_i m_j}\chi_{ij}/(2m_Z)$, with χ_{ij} a number of the order of unity at most. We are compelled to make this assumption due to our ignorance of the parameters involved in the scalar sector of the model. Although this can be a very optimistic assumption that can lead us to overestimate the scalar contribution to the $t \to c\gamma$ decay, one must have in mind that there is a suppression factor χ_{ct}, whose value could be very suppressed. In Fig. 4 we show the partial contribution from the neutral scalar boson accompanied by the c and t quarks to the C_R^ϕ coefficient as a function of m_ϕ. In this case $C_L^\phi = C_R^\phi$ and we do not show the contribution of the u quark as it is several orders of magnitude below than the c quark contribution. From this plot, we can also conclude that the scalar contribution to the $t \to c\gamma$

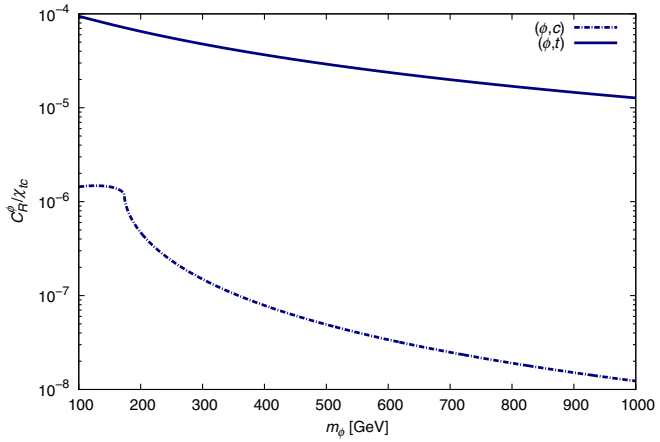

FIG. 4 (color online). Partial contribution from the loops carrying a neutral scalar boson and the c or t quarks to C_R^ϕ as a function of the scalar boson mass. For the $\phi\bar{c}t$ coupling we used the Cheng-Sher parametrization.

decay could be of the same order of magnitude than the gauge boson contribution. We also note that C_R^ϕ decreases rapidly as the scalar boson mass increases, but it depends considerably on the mass of the internal quark, which in fact is due to the use of the Cheng-Sher parametrization. It is also worth noting that the plateau observed in the case of the c quark contribution to C_R^ϕ is a reflection of the fact that below the mass threshold $m_\phi = m_t - m_c$, namely, $m_\phi \lesssim 174\,\mathrm{GeV}$, the t quark can decay as $t \to c\phi$, and so the Higgs-mediated $t \to c\gamma$ transition amplitude gets enhanced. Beyond this mass threshold, the $t \to c\phi$ decay is no longer kinematically allowed and C_R^ϕ becomes more suppressed as m_ϕ becomes heavier.

The individual contributions from a neutral scalar boson and the c and t quarks to $\mathrm{BR}(t \to c\gamma)$ are shown Fig. 5 as

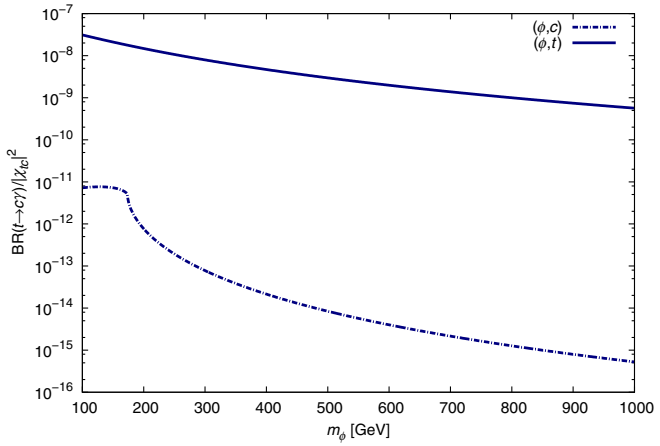

FIG. 5 (color online). Partial contributions from the loops carrying a neutral scalar boson ϕ and the c or t quarks to $\mathrm{BR}(t \to c\gamma)$ as a function of the scalar boson mass. For the $\phi\bar{c}t$ coupling we used the Cheng-Sher parametrization.

functions of the scalar boson mass. We observe that the t quark contribution is much larger than that of the c quark; hence, we expect that the bulk of the scalar contribution to $\mathrm{BR}(t \to c\gamma)$ would arise mainly from the loop carrying an internal top quark. Therefore the scalar contribution would be of the same order of magnitude as the gauge boson contribution.

V. FINAL REMARKS

We have analyzed the one-loop-induced decay $t \to c\gamma$ in the framework of 331 models, with particular emphasis on the minimal version. The generic contribution of these classes of models to this decay is induced by a new neutral gauge boson and a neutral scalar boson. In the minimal model there are also the contributions of singly and doubly charged gauge and scalar bosons accompanied by exotic quarks. We have found that, given the current constraints on the masses of the new particles, the dominant contribution to the $t \to c\gamma$ branching ratio could arise from the new charged gauge bosons and the lightest neutral scalar boson, although a branching ratio enhancement could be expected if the remaining scalars are also relatively light and have flavor-changing couplings. Contrary to the case of the contribution of the W gauge boson, the bilepton gauge boson contribution is free from large cancellations even if the bileptons are mass degenerate: there is an imperfect Glashow-Iliopoulos-Maiani-like mechanism in the minimal 331 model, which stems from the fact that the exotic quarks do not share the same electric charge. We examined a scenario in which $\mathrm{BR}(t \to c\gamma)$ could be of the order of 10^{-7}, but this value could be strongly suppressed as it has a large dependence on the values of the mixing matrix U_L that rotates up quarks from the flavor to the mass basis. For instance, if the U_L matrix elements are of the order of 10^{-1}, the bilepton contribution to the $t \to c\gamma$ branching ratio would be of the order 10^{-11}. In order to have an estimate for the contribution of the neutral scalar boson, we considered the Cheng-Sher ansatz for the flavor-changing couplings of the Higgs boson and found that the contribution to the $t \to c\gamma$ branching ratio could be of the order of 10^{-7} for a Higgs boson with a mass of the order of 100–200 GeV. A point worth mentioning is that, in 331 models without exotic quarks, the main contribution could arise from the lightest neutral scalar since the Z' mass is strongly constrained and so this contribution would be of the order of 10^{-9} at most.

As long as a particular 331 model was realized in nature, a more reliable estimate of the $t \to c\gamma$ decay would be obtained once more details of the model were known. We must conclude that any potential effects of 331 models on the $t \to c\gamma$ decay would hardly be observed in the near future.

ACKNOWLEDGMENTS

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TABLE I. Coupling constants for the flavor-changing vertices involving gauge bosons in the minimal 331 model. The right-handed coupling constants vanish and U_L stands for the mixing matrix that diagonalizes the SM up quarks.

Vertex $(\bar{q}_i q_j V)$	L_V^{ij}
$\bar{u}cZ'$	$\frac{c_W^2 U_{L31}^* U_{L32}}{\sqrt{3}\sqrt{1-4s_W^2}}$
$\bar{u}tZ'$	$\frac{c_W^2 U_{L31}^* U_{L33}}{\sqrt{3}\sqrt{1-4s_W^2}}$
$\bar{c}tZ'$	$\frac{c_W^2 U_{L32}^* U_{L33}}{\sqrt{3}\sqrt{1-4s_W^2}}$
$\bar{c}TY^-$	$\frac{c_W U_{L23}^*}{\sqrt{2}}$
$\bar{t}TY^-$	$\frac{c_W U_{L33}^*}{\sqrt{2}}$
$\bar{D}cY^{--}$	$\frac{c_W U_{L21}}{\sqrt{2}}$
$\bar{S}cY^{--}$	$\frac{c_W U_{L22}}{\sqrt{2}}$
$\bar{D}tY^{--}$	$\frac{c_W U_{L31}}{\sqrt{2}}$
$\bar{S}tY^{--}$	$\frac{c_W U_{L32}}{\sqrt{2}}$

APPENDIX: FEYNMAN RULES FOR THE $t \to c\gamma$ DECAY IN THE MINIMAL 331 MODEL

We first present the L_V^{ij} and R_V^{ij} coefficients necessary for the numerical evaluation of the $C_{L,R}^V$ coefficients in the minimal 331 model. The flavor conserving couplings of the Z' gauge boson to SM up quarks have the form of Eq. (15). The coefficients necessary for the calculation of the $t \to c\gamma$ amplitude can be extracted from Eq. (13) and are given by

$$L_{Z'}^{uu} = L_{Z'}^{cc} = -\frac{1-2s_W^2}{2\sqrt{3}\sqrt{1-4s_W^2}}, \tag{A1}$$

$$L_{Z'}^{tt} = \frac{1}{2\sqrt{3}\sqrt{1-4s_W^2}}, \tag{A2}$$

$$R_{Z'}^{uu} = R_{Z'}^{cc} = R_{Z'}^{tt} = \frac{2s_W^2}{\sqrt{3}\sqrt{1-4s_W^2}}. \tag{A3}$$

TABLE II. Feynman rules for the electromagnetic vertices involving charged gauge and scalar bosons. Q_Y (Q_ϕ) stands for the charge of the gauge (scalar) boson, and $\Gamma_{\alpha\beta\mu}(k_1,k_2,k_3) = (k_1 - k_2)_\mu g_{\alpha\beta} + (k_2 - k_3)_\alpha g_{\beta\mu} + (k_3 - k_1)_\beta g_{\mu\alpha}$, with all the momenta incoming.

Vertex	Feynman rule
$A_\mu(k_3)Y_\alpha(k_1)Y_\beta^\dagger(k_2)$	$-ieQ_Y\Gamma_{\alpha\beta\mu}(k_1,k_2,k_3)$
$Z'_\mu(k_3)Y_\alpha(k_1)Y_\beta^\dagger(k_2)$	$\frac{ig}{2c_W}\sqrt{3}\sqrt{1-4s_W^2}\Gamma_{\alpha\beta\mu}(k_1,k_2,k_3)$
$A_\mu\phi\phi^\dagger$	$-iQ_\phi(k_1-k_2)_\mu$

On the other hand, the coupling constants for the flavor-changing interactions of the heavy gauge bosons are purely left handed, as shown in Eqs. (16) and (17). The corresponding coupling constants are presented in Table I. Notice that the singly charged bilepton does not couple to a pair of SM quarks but only to a SM quark and an exotic quark.

For our calculation we also need the interaction of the photon with charged particles. Apart from the usual coupling of a photon with a fermion pair, $-ieQ_f\gamma^\mu$, the couplings of a charged gauge boson with the photon can be extracted from Eq. (12) and can be written as

$$\mathcal{L}^{YYA} = ieQ_Y(A^\mu(Y_{\mu\nu}Y^{\dagger\nu} - Y_{\mu\nu}^\dagger Y^\nu) - F_{\mu\nu}Y^\mu Y^{\dagger\nu}), \tag{A4}$$

with $Y_{\mu\nu} = \partial_\mu Y_\nu - \partial_\nu Y_\mu$. A similar term determines the interaction of the Z' gauge boson with the bilepton gauge bosons. As for the couplings of the photon with the new physical charged scalar bosons, they emerge from the kinetic term of the scalar triplets and sextet after rotating to the mass eigenstates. For a typical charged scalar boson ϕ we have

$$\mathcal{L}^{\phi\phi A} = (D_\mu^e\phi)^\dagger(D^{e\mu}\phi), \tag{A5}$$

where $D^{e\mu} = \partial^\mu + ieQ_\phi A^\mu$ is the $U_{\rm em}(1)$ covariant derivative. The Feynman rules for these vertices are presented in Table II.

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