PHYSICAL REVIEW D **87,** 075009 (2013)

Radiative seesaw-type mechanism of quark masses in $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$

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(Received 31 January 2013; published 12 April 2013)

We take up again the study of the mass spectrum of the quark sector in a model with gauge symmetry $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ (331). In a special type II-like 3-3-1 model, we obtain specific zero-texture mass matrices for the quarks which predict four massless quarks (u, c, d, s) and two massive quarks (b, t) at the electroweak scale ($\sim$ GeV). By considering the mixing between the standard model quarks and new exotic quarks at large scales predicted by the model, we find that a third quark (associated to the charm quark) acquires a mass. The remaining light quarks (u, d, s) get small masses ($\sim$ MeV) via radiative corrections.

DOI: 10.1103/PhysRevD.87.075009 PACS numbers: 12.60.Cn, 11.15.Ex, 11.30.Qc, 12.15.Ff

I. INTRODUCTION

The ATLAS and CMS experiments at the CERN Large Hadron Collider (LHC) have found a 126 GeV Higgs boson [1–4] through the $h \to \gamma\gamma$ decay channel, increasing our knowledge of the electroweak symmetry breaking (EWSB) sector and opening a new era in particle physics. Now the priority of the LHC experiments will be to measure precisely the couplings of the new particle to standard model (SM) fermions and gauge bosons and to establish its quantum numbers. It also remains to look for further new states associated with the EWSB mechanism which will allow us to discriminate among the different theoretical models addressed to explain EWSB.

Despite all its success, the SM of the electroweak interactions based on the $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry has many unexplained features [5]. Most of them are linked to the mechanism responsible for the stabilization of the weak scale, the origin of fermion masses and mixings, and the three family structure. Because of this reason, many people consider the standard model to be an effective framework of a yet unknown more fundamental theory. A fundamental theory, one expects, should have a dynamical explanation for the masses and mixings. The lack of predictivity of the fermion masses and mixings in the SM has motivated many models based on extended symmetries in the context of two Higgs doublets, grand unification, extra dimensions, and superstrings leading to specific textures for the Yukawa couplings [6–8]. The understanding of the discrete flavor symmetries hidden in such textures may be useful in understanding the underlying dynamics responsible for quark mass generation and *CP* violation. One clear and outstanding feature in the pattern of quark masses is that they increase from one generation to the next spreading over a range of 5 orders of magnitude, and the mixings from the first to the second and to the third family are in decreasing order [9–12]. From the phenomenological point of view, it is possible to describe some features of the mass hierarchy by assuming zero-texture Yukawa matrices [13]. Models with spontaneously broken flavor symmetries may also produce hierarchical mass structures. These horizontal symmetries can be continuous and Abelian, as the original Froggatt-Nielsen model [14], or non-Abelian as, for example, $SU(3)$ and $SO(3)$ family models [15]. Models with discrete symmetries may also predict mass hierarchies for leptons [16] and quarks [17]. Other models with horizontal symmetries have been proposed in the literature [18].

On the other hand, the origin of the family structure of the fermions can be addressed in family-dependent models where a symmetry distinguishes fermions of different families. Alternatively, an explanation for this issue can also be provided by the models based on the gauge symmetry $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$, also called 3-3-1 models, which introduce a family nonuniversal $U(1)_X$ symmetry [19–22]. These models have a number of phenomenological advantages. First of all, the three family structure in the fermion sector can be understood in the 3-3-1 models from the cancellation of chiral anomalies [23] and asymptotic freedom in QCD. Secondly, the fact that the third family is treated under a different representation can explain the large mass difference between the heaviest quark family and the two lighter ones [24]. Finally, these models contain a natural Peccei-Quinn symmetry that is necessary to solve the strong-*CP* problem [25].

The 3-3-1 models extend the scalar sector of the SM into three $SU(3)_L$ scalar triplets: one heavy triplet field with a vacuum expectation value (VEV) at high energy scale $\langle\chi\rangle = \nu_\chi$, which breaks the symmetry $SU(3)_L \otimes U(1)_X$ into the SM electroweak group $SU(2)_L \otimes U(1)_Y$,

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and two lighter triplets with VEVs at the electroweak scale $\langle\rho\rangle = v_\rho$ and $\langle\eta\rangle = v_\eta$, which trigger electroweak symmetry breaking. Besides that, the 3-3-1 model could possibly explain the excess of events in the $h \to \gamma\gamma$ decay recently observed at the LHC, since the heavy exotic quarks, the charged Higgs, and the heavy charged gauge bosons contribute to this process. On the other hand, the 3-3-1 model reproduces a specialized two Higgs doublet model type III (2HDM-III) in the low energy limit, where both triplets ρ and η are decomposed into two hypercharge-one $SU(2)_L$ doublets plus charged and neutral singlets. Thus, like the 2HDM-III, the 3-3-1 model can predict huge flavor changing neutral currents and CP-violating effects, which are severely suppressed by experimental data at electroweak scales. In the 2HDM-III, for each quark type, up or down, there are two Yukawa couplings. One of the Yukawa couplings is for generating the quark masses, and the other one produces the flavor changing couplings at tree level. One way to remove both the huge flavor changing neutral currents and CP-violating effects is by imposing discrete symmetries obtaining two types of 3-3-1 models (type I and II models) which exhibit the same Yukawa interactions as the 2HDM type I and II at low energy where each fermion is coupled at most to one Higgs doublet. In the 3-3-1 model type I, one Higgs electroweak triplet (for example, ρ) provides masses to the phenomenological up- and down-type quarks, simultaneously. In the type II, one Higgs triplet (η) gives masses to the up-type quarks and the other triplet (ρ) to the down-type quarks. Recently, the authors in Ref. [26] discussed the mass structures in the framework of the I-type 3-3-1 model. In this paper, we obtain different structures for the type II-like model. We found that only the top and bottom quarks acquire masses if the mixing of the SM quarks with the exotic quarks is neglected. We obtain by the method of recursive expansion [27] that if mixing couplings with the heavy quark sector of the 3-3-1 model are considered, only the charm quark obtains a mass, while the light quarks remain massless. The masses of the up, down, and strange quarks are generated through loop corrections which are kind of seesaw-type radiative mechanisms that involve the virtual exotic quarks as well as neutral and charged scalars running in the loops. Thus, the hierarchy of the quark mass spectrum can be explained from three different sources: the tree-level quark mass matrices from the symmetry breaking, the mixings between the SM quarks and the exotic quarks, and seesaw-type radiative corrections. This mechanism of generating the quark masses provides an alternative to the ones discussed in Refs. [28,29] where effective operators and one-loop corrections are introduced.

This paper is organized as follows. In Sec. II, we briefly describe some theoretical aspects of the 3-3-1 model and its particle content, in particular, in the fermionic and scalar sector in order to obtain the mass spectrum. Section III discusses the possible zero textures for the SM quark mass matrices at tree level. In Sec. IV, we obtain the quark masses at tree and one-loop level of the complete model by imposing specific zero-texture masses. Finally in Sec. V, we state our conclusions.

II. THE FERMION AND SCALAR SECTOR

We consider the 3-3-1 model where the electric charge is defined by

$$Q = T_3 - \frac{1}{\sqrt{3}}T_8 + X, \tag{1}$$

with $T_3 = \frac{1}{2}\mathrm{diag}(1,-1,0)$ and $T_8 = (\frac{1}{2\sqrt{3}})\mathrm{diag}(1,1,-2)$. In order to avoid chiral anomalies, the model introduces in the fermionic sector the following $[SU(3)_c, SU(3)_L, U(1)_X]$ left-handed representations:

$$\begin{aligned}
Q_L^1 &= \begin{pmatrix} U^1 \\ D^1 \\ T^1 \end{pmatrix}_L : (3,3,1/3), \quad \begin{cases} U_R^1: (3^*,1,2/3) \\ D_R^1: (3^*,1,-1/3) \\ T_R^1: (3^*,1,2/3), \end{cases} \\
Q_L^{2,3} &= \begin{pmatrix} D^{2,3} \\ U^{2,3} \\ J^{2,3} \end{pmatrix}_L : (3,3^*,0), \quad \begin{cases} D_R^{2,3}: (3^*,-1/3) \\ U_R^{2,3}: (3^*,1,2/3) \\ J_R^{2,3}: (3^*,1,-1/3), \end{cases} \\
L_L^{1,2,3} &= \begin{pmatrix} \nu^{1,2,3} \\ e^{1,2,3} \\ (\nu^{1,2,3})^c \end{pmatrix}_L : (1,3,-1/3), \quad \begin{cases} e_R^{1,2,3}: (1,1,-1) \\ N_R^{1,2,3}: (1,1,0), \end{cases}
\end{aligned} \tag{2}$$

where U_L^i and D_L^i for $i = 1, 2, 3$ are three up- and down-type quark components in the flavor basis, while ν_L^i and e_L^i are the neutral and charged lepton families. The right-handed sector transforms as singlets under $SU(3)_L$ with $U(1)_X$ quantum numbers equal to the electric charges. In addition, we see that the model introduces heavy fermions with the following properties: a single flavor quark T^1 with electric charge $2/3$, two flavor quarks $J^{2,3}$ with charge $-1/3$, three neutral Majorana leptons $(\nu^{1,2,3})_L^c$, and three right-handed Majorana leptons $N_R^{1,2,3}$. On the other hand, the scalar sector introduces one triplet field with VEV $\langle\chi\rangle_0 = v_\chi$, which provides the masses to the new heavy fermions, and two triplets with VEVs $\langle\rho\rangle_0 = v_\rho$ and $\langle\eta\rangle_0 = v_\eta$, which give masses to the SM fermions at the electroweak scale. However, as it will be shown in Sec. IV, we can have a discrete symmetry in the quark sector that allows the triplet χ to give masses not only to the heavy exotic quarks but also to the light quarks via a radiative seesaw-type mechanism while the triplets ρ and η give masses to the remaining quarks. The $[SU(3)_L, U(1)_X]$ group structure of the scalar fields is

$$\chi = \begin{pmatrix} \chi_1^0 \\ \chi_2^- \\ \frac{1}{\sqrt{2}}(\upsilon_\chi + \xi_\chi \pm i\zeta_\chi) \end{pmatrix}: (3, -1/3), \qquad \rho = \begin{pmatrix} \rho_1^+ \\ \frac{1}{\sqrt{2}}(\upsilon_\rho + \xi_\rho \pm i\zeta_\rho) \\ \rho_3^+ \end{pmatrix}: (3, 2/3),$$
$$\eta = \begin{pmatrix} \frac{1}{\sqrt{2}}(\upsilon_\eta + \xi_\eta \pm i\zeta_\eta) \\ \eta_2^- \\ \eta_3^0 \end{pmatrix}: (3, -1/3). \tag{3}$$

The EWSB follows the scheme $SU(3)_L \otimes U(1)_X \xrightarrow{\langle\chi\rangle} SU(2)_L \otimes U(1)_Y \xrightarrow{\langle\eta\rangle,\langle\rho\rangle} U(1)_Q$, where the vacuum expectation values satisfy $\upsilon_\chi \gg \upsilon_\eta, \upsilon_\rho$.

The interactions among the scalar fields are contained in the following most general potential that we can construct with three scalar triplets:

$$\begin{aligned} V_H = {} & \mu_\chi^2(\chi^\dagger\chi) + \mu_\eta^2(\eta^\dagger\eta) + \mu_\rho^2(\rho^\dagger\rho) + f(\chi_i\eta_j\rho_k\varepsilon^{ijk} + \text{H.c.}) + \lambda_1(\chi^\dagger\chi)(\chi^\dagger\chi) \\ & + \lambda_2(\rho^\dagger\rho)(\rho^\dagger\rho) + \lambda_3(\eta^\dagger\eta)(\eta^\dagger\eta) + \lambda_4(\chi^\dagger\chi)(\rho^\dagger\rho) + \lambda_5(\chi^\dagger\chi)(\eta^\dagger\eta) \\ & + \lambda_6(\rho^\dagger\rho)(\eta^\dagger\eta) + \lambda_7(\chi^\dagger\eta)(\eta^\dagger\chi) + \lambda_8(\chi^\dagger\rho)(\rho^\dagger\chi) + \lambda_9(\rho^\dagger\eta)(\eta^\dagger\rho). \end{aligned} \tag{4}$$

After the symmetry breaking, it is found that the mass eigenstates are related to the weak states in the scalar sector by [21,22]

$$\begin{pmatrix} G_1^\pm \\ H_1^\pm \end{pmatrix} = R_{\beta_T}\begin{pmatrix} \rho_1^\pm \\ \eta_2^\pm \end{pmatrix}, \qquad \begin{pmatrix} G_1^0 \\ A_1^0 \end{pmatrix} = R_{\beta_T}\begin{pmatrix} \zeta_\rho \\ \zeta_\eta \end{pmatrix},$$
$$\begin{pmatrix} H_1^0 \\ h^0 \end{pmatrix} = R_{\alpha_T}\begin{pmatrix} \xi_\rho \\ \xi_\eta \end{pmatrix}, \tag{5}$$

$$\begin{pmatrix} G_2^0 \\ H_2^0 \end{pmatrix} = R\begin{pmatrix} \chi_1^0 \\ \eta_3^0 \end{pmatrix}, \qquad \begin{pmatrix} G_3^0 \\ H_3^0 \end{pmatrix} = \frac{-1}{\sqrt{2}}R\begin{pmatrix} \zeta_\chi \\ \xi_\chi \end{pmatrix},$$
$$\begin{pmatrix} G_2^\pm \\ H_2^\pm \end{pmatrix} = R\begin{pmatrix} \chi_2^\pm \\ \rho_3^\pm \end{pmatrix}, \tag{6}$$

with

$$R_{\alpha_T(\beta_T)} = \begin{pmatrix} \cos\alpha_T(\beta_T) & \sin\alpha_T(\beta_T) \\ -\sin\alpha_T(\beta_T) & \cos\alpha_T(\beta_T) \end{pmatrix},$$
$$R = \begin{pmatrix} -1 & 0 \\ 0 & 1 \end{pmatrix}, \tag{7}$$

where $\tan\beta_T = \upsilon_\eta/\upsilon_\rho$, and $\tan 2\alpha_T = M_1/(M_2 - M_3)$ with

$$\begin{aligned} M_1 &= 4\lambda_6\upsilon_\eta\upsilon_\rho + 2\sqrt{2}f\upsilon_\chi, \\ M_2 &= 4\lambda_2\upsilon_\rho^2 - \sqrt{2}f\upsilon_\chi\tan\beta_T, \\ M_3 &= 4\lambda_3\upsilon_\eta^2 - \sqrt{2}f\upsilon_\chi/\tan\beta_T. \end{aligned} \tag{8}$$

With the above spectrum, we obtain the following $SU(3)_L \otimes U(1)_X$ renormalizable Yukawa Lagrangian for the quark sector [26]:

$$\begin{aligned} -\mathcal{L}_Y = {} & \bar{Q}_L^1(\eta h_{\eta 1j}^U + \chi h_{\chi 1j}^U)U_R^j + \bar{Q}_L^1\rho h_{\rho 1j}^D D_R^j \\ & + \bar{Q}_L^1\rho h_{\rho 1m}^J J_R^m + \bar{Q}_L^1(\eta h_{\eta 11}^T + \chi h_{\chi 11}^T)T_R^1 \\ & + \bar{Q}_L^n\rho^* h_{\rho nj}^U U_R^j + \bar{Q}_L^n(\eta^* h_{\eta nj}^D + \chi^* h_{\chi nj}^D)D_R^j \\ & + \bar{Q}_L^n(\eta^* h_{\eta nm}^J + \chi^* h_{\chi nm}^J)J_R^m + \bar{Q}_L^n\rho^* h_{\rho n1}^T T_R^1 + \text{H.c.}, \end{aligned} \tag{9}$$

where $n = 2, 3$ is the index that labels the second and third quark triplets shown in Eq. (2), and $h_{\phi ij}^f$ are the i, j components of nondiagonal matrices in the flavor space associated with each scalar triplet ϕ: η, ρ, χ.

III. ZERO-TEXTURE MASSES AT LOW ENERGY

By considering a scenario where the mixing terms among fields at small and large mass scales in Eq. (9) do not contribute at low energy, we obtain the following decoupled low energy Yukawa Lagrangian:

$$\begin{aligned} -\mathcal{L}_Y^{\text{LE}} = {} & \bar{Q}_L^1\eta h_{\eta 1j}^U U_R^j + \bar{Q}_L^n\eta^* h_{\eta nj}^D D_R^j + \bar{Q}_L^n\rho^* h_{\rho nj}^U U_R^j \\ & + \bar{Q}_L^1\rho h_{\rho 1j}^D D_R^j + \text{H.c.} \end{aligned}$$

Although the above Lagrangian exhibits the same general form as the 2HDM Lagrangian, they are not the same because the Abelian sector $U(1)_X$ of the 3-3-1 symmetry introduces a quantum number that differentiates the second and third rows from the first one, from which not all Yukawa couplings are allowed by the symmetry. In 2HDM models, there are no labels and the three rows are identical. Thus, by imposing appropriate discrete symmetries, many *Ansätze* for the couplings can be obtained. From the previous expression it follows that the mass Lagrangian corresponding to the SM quark sector is given by

$$-\mathcal{L}_Y^{\text{mass-SM}} = \frac{v_\eta}{\sqrt{2}}\bar{U}_L^1 h_{\eta 1j}^U U_R^j + \frac{v_\eta}{\sqrt{2}}\bar{D}_L^n h_{\eta nj}^D D_R^j + \frac{v_\rho}{\sqrt{2}}\bar{U}_L^n h_{\rho nj}^U U_R^j + \frac{v_\rho}{\sqrt{2}}\bar{D}_L^1 h_{\rho 1j}^D D_R^j + \text{H.c.} \tag{10}$$

The 3-3-1 model gives the different possible textures that can be chosen independently for the up and down sector according to the discrete symmetry imposed. These textures are given by

$$M_U^{(A)} = \frac{v_\eta}{\sqrt{2}}\begin{pmatrix} h_{\eta 11}^U & h_{\eta 1m}^U \\ 0_{2\times 1} & 0_{2\times 2}\end{pmatrix}, \quad M_U^{(B)} = \frac{v_\rho}{\sqrt{2}}\begin{pmatrix} 0 & 0_{1\times 2} \\ h_{\rho n1}^U & h_{\rho nm}^U\end{pmatrix}, \tag{11}$$

$$M_D^{(A)} = \frac{v_\eta}{\sqrt{2}}\begin{pmatrix} 0 & 0_{1\times 2} \\ h_{\eta n1}^D & h_{\eta nm}^D\end{pmatrix}, \quad M_D^{(B)} = \frac{v_\rho}{\sqrt{2}}\begin{pmatrix} h_{\rho 11}^D & h_{\rho 1m}^D \\ 0_{2\times 1} & 0_{2\times 2}\end{pmatrix}, \tag{12}$$

with n and $m = 2, 3$. The choice of the textures $M_U^{(A)}$ and $M_D^{(B)}$ (where the subindices U and D refer to the up and down sector, respectively) can be obtained by imposing

$$U_R \to -U_R, \qquad D_R \to D_R, \qquad \rho \to \rho, \qquad \eta \to -\eta. \tag{13}$$

These textures imply that the quarks generate mass from the Q_L^1 terms at tree level. In this case, only the top and bottom quarks would acquire mass while the other quarks will remain massless. The textures $M_U^{(B)}$ and $M_D^{(A)}$ can be obtained through

$$U_R \to -U_R, \qquad D_R \to D_R, \qquad \rho \to -\rho, \qquad \eta \to \eta. \tag{14}$$

These would imply four massive quarks from the Q_L^n terms at tree level, which will be unnatural since the masses will be practically generated by hand through the Yukawa couplings. The choice of $M_U^{(A)}$ and $M_D^{(A)}$, where

$$U_R \to -U_R, \qquad D_R \to -D_R, \qquad \rho \to \rho, \qquad \eta \to -\eta \tag{15}$$

will generate a mass for the top quark, which comes from the Q_L^1 term, while the bottom quark coming from the Q_L^n terms will be massless and the d and s quarks will be massive at tree level. Thus, the choice of this textures does not lead to a phenomenological viable quark mass spectrum. The choice of $M_U^{(B)}$ and $M_D^{(B)}$ with

$$U_R \to -U_R, \qquad D_R \to -D_R, \qquad \rho \to -\rho, \qquad \eta \to \eta \tag{16}$$

would imply that only three quarks will be massive. In the up sector, one can choose the massive top and charm quarks as elements of Q_L^n. In that case, the bottom and strange will be massless while the down quark coming from Q_L^1 will acquire mass, which is unnatural. In conclusion, the textures $M_U^{(A)}$ and $M_D^{(B)}$ with the discrete symmetry in Eq. (13) could provide a better explanation for the quark mass hierarchy. The vanishing entries will be filled by the mixings between the SM and exotic quarks or by radiative corrections.

IV. ZERO-TEXTURE MASSES WITH MIXING COUPLINGS

In order to obtain the submatrices $M_U^{(A)}$ and $M_D^{(B)}$ in Eqs. (11) and (12) from the original 3-3-1 Lagrangian in (9), we extend the discrete symmetry in (13) to

$$U_R \to -U_R, \qquad D_R \to D_R, \qquad \eta \to -\eta, \qquad \rho \to \rho, \qquad \chi \to \chi, \qquad T_R \to T_R, \qquad J_R \to J_R, \tag{17}$$

which restricts the $SU(3)_L \times U(1)_X$ renormalizable Yukawa Lagrangian to be given by

$$-\mathcal{L}_Y = \overline{Q}_L^1 \eta h_{\eta 1j}^U U_R^j + \overline{Q}_L^1 \rho h_{\rho 1j}^D D_R^j + \overline{Q}_L^1 \chi h_{\chi 11}^T T_R^1 + \overline{Q}_L^n \chi^* h_{\chi nm}^J J_R^m + \overline{Q}_L^n \rho^* h_{\rho n1}^T T_R^1 + \overline{Q}_L^n \chi^* h_{\chi nj}^D D_R^j + \overline{Q}_L^1 \rho h_{\rho 1m}^J J_R^m + \text{H.c.} \tag{18}$$

The previous Lagrangian can be rewritten as

$$-\mathcal{L}_Y = -\mathcal{L}_Y^{(1)} - \mathcal{L}_Y^{(2)} - \mathcal{L}_Y^{(3)}, \tag{19}$$

where $-\mathcal{L}_Y^{(1)}$ corresponds to the quark mass Lagrangian, while $-\mathcal{L}_Y^{(2)}$ and $-\mathcal{L}_Y^{(3)}$ are the Lagrangians which include the interactions of the quarks with the neutral and charged Higgs and Goldstone bosons, respectively. These Lagrangians are given by

$$-\mathcal{L}_Y^{(1)} = \frac{v_\eta}{\sqrt{2}}\overline{U}_L^1 h_{\eta 1j}^U U_R^j + \frac{v_\rho}{\sqrt{2}}\overline{U}_L^n h_{\rho n1}^T T_R^1 + \frac{v_\chi}{\sqrt{2}}\overline{T}_L^1 h_{\chi 11}^T T_R^1 + \frac{v_\rho}{\sqrt{2}}\overline{D}_L^1 h_{\rho 1j}^D D_R^j + \frac{v_\rho}{\sqrt{2}}\overline{D}_L^1 h_{\rho 1m}^J J_R^m + \frac{v_\chi}{\sqrt{2}}\overline{J}_L^n h_{\chi nj}^D D_R^j + \frac{v_\chi}{\sqrt{2}}\overline{J}_L^n h_{\chi nm}^J J_R^m + \text{H.c.}, \tag{20}$$

$$-\mathcal{L}_Y^{(2)} = \frac{1}{\sqrt{2}}\overline{U}_L^1(\xi_\eta + i\zeta_\eta)h^U_{\eta 1j}U_R^j + \overline{U}_L^1\chi_1^0 h^T_{\chi 11}T_R^1 + \overline{T}_L^1\eta_3^0 h^U_{\eta 1j}U_R^j + \frac{1}{\sqrt{2}}\overline{T}_L^1(\xi_\chi + i\zeta_\chi)h^T_{\chi 11}T_R^1 + \frac{1}{\sqrt{2}}\overline{U}_L^n(\xi_\rho - i\zeta_\rho)h^T_{\rho n1}T_R^1 + \frac{1}{\sqrt{2}}\overline{D}_L^1(\xi_\rho + i\zeta_\rho)h^D_{\rho 1j}D_R^j + \frac{1}{\sqrt{2}}\overline{D}_L^1(\xi_\rho + i\zeta_\rho)h^J_{\rho 1m}J_R^m + \overline{D}_L^n\chi_1^0 h^D_{\chi nj}D_R^j + \overline{D}_L^n\chi_1^0 h^J_{\chi nm}J_R^m + \frac{1}{\sqrt{2}}\overline{J}_L^n(\xi_\chi - i\zeta_\chi)h^D_{\chi nj}D_R^j + \frac{1}{\sqrt{2}}\overline{J}_L^n(\xi_\chi - i\zeta_\chi)h^J_{\chi nm}J_R^m + \text{H.c.}, \tag{21}$$

$$-\mathcal{L}_Y^{(3)} = +\overline{D}_L^n\rho_1^- h^T_{\rho n1}T_R^1 + \overline{J}_L^n\rho_3^- h^T_{\rho n1}T_R^1 + \overline{U}_L^1\rho_1^+ h^D_{\rho 1j}D_R^j + \overline{T}_L^1\rho_3^+ h^D_{\rho 1j}D_R^j + \overline{U}_L^n\chi_2^+ h^D_{\chi nj}D_R^j + \overline{U}_L^1\rho_1^+ h^J_{\rho 1m}J_R^m + \overline{T}_L^1\rho_3^+ h^J_{\rho 1m}J_R^m + \overline{U}_L^n\chi_2^+ h^J_{\chi nm}J_R^m + \text{H.c.} \tag{22}$$

From Eq. (20) it follows that the mass matrices for the up- and down-type quarks are given by

$$M^U = \left(\begin{array}{ccc|c} v_\eta h^U_{\eta 11} & v_\eta h^U_{\eta 12} & v_\eta h^U_{\eta 13} & 0 \\ 0 & 0 & 0 & v_\rho h^T_{\rho 21} \\ 0 & 0 & 0 & v_\rho h^T_{\rho 31} \\ \hline 0 & 0 & 0 & v_\chi h^T_{\chi 11} \end{array}\right) = \left(\begin{array}{c|c} M_U^{(A)} & k_{3\times 1} \\ \hline 0_{1\times 3} & M_T \end{array}\right),$$

$$M^D = \left(\begin{array}{ccc|cc} v_\rho h^D_{\rho 11} & v_\rho h^D_{\rho 12} & v_\rho h^D_{\rho 13} & v_\rho h^J_{\rho 12} & v_\rho h^J_{\rho 13} \\ 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \\ \hline v_\chi h^D_{\chi 21} & v_\chi h^D_{\chi 22} & v_\chi h^D_{\chi 23} & v_\chi h^J_{\chi 22} & v_\chi h^J_{\chi 23} \\ v_\chi h^D_{\chi 31} & v_\chi h^D_{\chi 32} & v_\chi h^D_{\chi 33} & v_\chi h^J_{\chi 32} & v_\chi h^J_{\chi 33} \end{array}\right) = \left(\begin{array}{c|c} M_D^{(B)} & s_{3\times 2} \\ \hline S_{2\times 3} & M_J \end{array}\right), \tag{23}$$

where the diagonal blocks $M_U^{(A)}$ and $M_D^{(B)}$ are the same as (11) and (12), $M_{T,J}$ are the masses of the T^1 and J^n quarks, and k, s, and S are mixing mass blocks. The different VEVs of the scalars have the following hierarchy:

$$v_\chi >> v_\rho, \qquad v_\eta \sim 246\ \text{GeV}. \tag{24}$$

A. Up sector

The mass matrix for the up-type quarks satisfies the following relation:

$$M^U(M^U)^\dagger = \begin{pmatrix} m_t^2 & 0_{1\times 3} \\ 0_{3\times 1} & \tilde{M} \end{pmatrix}, \qquad m_t^2 = v_\eta^2\sum_{i=1}^{3}|h^U_{\eta 1i}|^2, \tag{25}$$

where $\tilde{M}$ is given by

$$\tilde{M} = \left(\begin{array}{cc|c} v_\rho^2|h^T_{\rho 21}|^2 & v_\rho^2 h^T_{\rho 21}(h^T_{\rho 31})^* & v_\rho v_\chi h^T_{\rho 21}(h^T_{\chi 11})^* \\ v_\rho^2 h^T_{\rho 31}(h^T_{\rho 21})^* & v_\rho^2|h^T_{\rho 31}|^2 & v_\rho v_\chi h^T_{\rho 31}(h^T_{\chi 11})^* \\ \hline v_\rho v_\chi h^T_{\chi 11}(h^T_{\rho 21})^* & v_\rho v_\chi h^T_{\chi 11}(h^T_{\rho 31})^* & v_\chi^2|h^T_{\chi 11}|^2 \end{array}\right) = \left(\begin{array}{c|c} h^T_{\rho n1}(h^T_{\rho m1})^\dagger v_\rho^2 & h^T_{\rho n1}(h^T_{\chi 11})^* v_\rho v_\chi \\ \hline h^T_{\chi 11}(h^T_{\rho m1})^\dagger v_\rho v_\chi & |h^T_{\chi 11}|^2 v_\chi^2 \end{array}\right). \tag{26}$$

The submatrix $\tilde{M}$ satisfies the following relation:

$$\det(\tilde{M}) = 0. \tag{27}$$

Therefore, one quark (the u quark) remains massless at tree level; one quark (the c quark) will acquire mass through the mixing with the exotic T quark, and two quarks (the t and T quarks) have tree-level masses without mixing. To find the rotation matrix which diagonalizes the matrix $\tilde{M}$, we perform a perturbative diagonalization. The mass matrix $\tilde{M}$ can be block diagonalized through the rotation matrix W_L according to

$$W_L^\dagger \tilde{M} W_L \simeq \begin{pmatrix} \tilde{f} & 0_{2\times 1} \\ 0_{1\times 2} & m_T^2 \end{pmatrix}, \tag{28}$$

with

$$W_L = \begin{pmatrix} 1_{2\times 2} & B \\ -B^\dagger & 1 \end{pmatrix}, \qquad m_T^2 \simeq |h^T_{\chi 11}|^2 v_\chi^2, \tag{29}$$

where $\tilde{f}_{nm} = -v_\rho^2 h^T_{\rho n1}(h^T_{\rho m1})^\dagger$ with $m, n = 2, 3$. From the condition of the vanishing of the off-diagonal submatrices in the previous expression, we obtain at leading order in B the following relations:

$$aB + b - Bm_T^2 = 0, \qquad B^\dagger a + b^\dagger - m_T^2 B^\dagger = 0,$$

where a and b have the following components:

$$a_{nm} = h^T_{\rho n1}(h^T_{\rho m1})^\dagger v_\rho^2, \qquad b_{n1} = h^T_{\rho n1}(h^T_{\chi 11})^* v_\rho v_\chi. \tag{30}$$

By using the method of recursive expansion, taking into account the hierarchy $a_{nm} << b_{n1} << m_T^2$, we find that the submatrix B is approximately given by

$$B_{n1} \simeq \frac{v_\rho}{v_\chi}\frac{h^T_{\rho n1}}{h^T_{\chi 11}} \simeq \frac{m_c}{m_T}. \tag{31}$$

For the sake of simplicity, let us assume that the Yukawa couplings $h^T_{\rho m1}$ are real. In that case, the matrix $\tilde{f}$ is diagonalized by a rotation matrix

$$V_{Luc} = \frac{1}{\sqrt{(h^T_{\rho 11})^2 + (h^T_{\rho 21})}}\begin{pmatrix} h^T_{\rho 21} & -h^T_{\rho 31} \\ h^T_{\rho 31} & h^T_{\rho 11} \end{pmatrix} \tag{32}$$

according to

$$V^T_{Luc}\tilde{f}V_{Luc} = \tilde{f}_{\rm diag} = {\rm diag}(-m_c^2, 0), \qquad m_c^2 = v_\rho^2[(h^T_{\rho 21})^2 + (h^T_{\rho 31})^2], \qquad m_u = 0. \tag{33}$$

Since $v_\rho \sim v_\eta$, it follows that $|h^U_{\eta 1i}| >> |h^T_{\rho m1}|$. Here the following identity has been taken into account:

$$\begin{pmatrix} \frac{c}{\sqrt{c^2+d^2}} & \frac{d}{\sqrt{c^2+d^2}} \\ -\frac{d}{\sqrt{c^2+d^2}} & \frac{c}{\sqrt{c^2+d^2}} \end{pmatrix}\begin{pmatrix} c^2 & cd \\ cd & d^2 \end{pmatrix}\begin{pmatrix} \frac{c}{\sqrt{c^2+d^2}} & -\frac{d}{\sqrt{c^2+d^2}} \\ \frac{d}{\sqrt{c^2+d^2}} & \frac{c}{\sqrt{c^2+d^2}} \end{pmatrix} = \begin{pmatrix} c^2+d^2 & 0 \\ 0 & 0 \end{pmatrix}. \tag{34}$$

Then, it follows that the mass matrix $\tilde{M}$ is diagonalized by a rotation matrix R_L according to

$$R_L^T\tilde{M}R_L = {\rm diag}(-m_c^2, 0, m_T^2), \quad \text{with} \quad R_L = \begin{pmatrix} V_{Luc} & B \\ -B^\dagger V_{Luc} & 1 \end{pmatrix}. \tag{35}$$

Therefore, the mass matrix $M^U(M^U)^\dagger$ is diagonalized by a rotation matrix V_L^U according to

$$(V_L^U)^\dagger M^U(M^U)^\dagger V_L^U = {\rm diag}(m_t^2, -m_c^2, 0, m_T^2), \quad \text{with} \quad V_L^U = \begin{pmatrix} 1 & 0_{1\times 3} \\ 0_{3\times 1} & R_L \end{pmatrix}. \tag{36}$$

Although the mixing terms with the exotic sector allow nonvanishing mass to the c quark, the lightest quark u remains massless due to the zero texture of M^U, as shown in Eq. (23). However, the vanishing entries can be filled by radiative corrections. For the sake of simplicity, we assume that the CP-odd neutral scalars are much heavier than the heavy exotic quarks T and J^2, so that their loop contributions to the entries of the quark mass matrix can be neglected. Here we do not consider the contributions coming from the exotic quark J^3 since we assume that it does not mix with the SM quarks and with the exotic quark J^2. Then, the heavy exotic quark T with the neutral scalars ξ_ρ, ξ_χ, η_3^0, and the heavy exotic quark J^2 with the charged scalars ρ_1^+ and $\rho_3^\pm$ running in the loop induce radiative corrections at one-loop level to most of the entries of the up-type quark mass matrix, thanks to the scalar quartic interactions. These virtual scalars couple to real neutral scalars which acquire VEVs after electroweak symmetry breaking. In this manner, the up quark mass is radiatively generated in an analogous way to the loop induced neutrino mass generation processes. Besides that, we assume that the quartic scalar couplings are approximately equal. Here we also assume that $h^T_{\chi 11} >> |h^U_{\eta 1i}|$, $|h^T_{\rho m1}|$, and $h^J_{\chi mm}(m = 2, 3)$ is much bigger than the magnitudes of the remaining down-type quark Yukawa couplings. These assumptions allow us to neglect the loop contributions to the up- and down-type quark mass matrices that involve the mixings between the SM quarks and the exotic quarks in the internal lines. Therefore, the leading one-loop level contributions to the entries of the up-type quark mass matrix come from the Feynman diagrams shown in Fig. 1. Here we use the unitary gauge where we get rid of the Goldstone bosons $G_1^\pm$, $G_2^\pm$, G_1^0, G_2^0, and G_3^0. Hence, the radiative corrections constraint the up-type quark mass matrix to be of the form

$$M^U \begin{pmatrix} v_\eta h^U_{\eta 11} & v_\eta h^U_{\eta 12} & v_\eta h^U_{\eta 13} & (\delta M^U)_{14} \\ (\delta M^U)_{21} & (\delta M^U)_{22} & (\delta M^U)_{23} & v_\rho h^T_{\rho 21} + (\delta M^U)_{24} \\ (\delta M^U)_{31} & (\delta M^U)_{32} & (\delta M^U)_{33} & v_\rho h^T_{\rho 31} + (\delta M^U)_{34} \\ (\delta M^U)_{41} & (\delta M^U)_{42} & (\delta M^U)_{43} & v_\chi h^T_{\chi 11} + (\delta M^U)_{44} \end{pmatrix}, \tag{37}$$

where their dominant loop-induced entries are given by

$$(\delta M^U)_{14} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^J_{\rho 12}h^T_{\rho 21}v_\chi^2}{m_{J_2}}C_0\left(\frac{m_{\rho_1^+}}{m_{J_2}}, \frac{m_{\rho_3^-}}{m_{J_2}}\right),$$
$$(\delta M^U)_{m1} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1}h^U_{\eta 11}v_\eta v_\rho}{m_T}C_0\left(\frac{m_{\xi_\rho}}{m_T}, \frac{m_{\eta_3^0}}{m_T}\right), \tag{38}$$

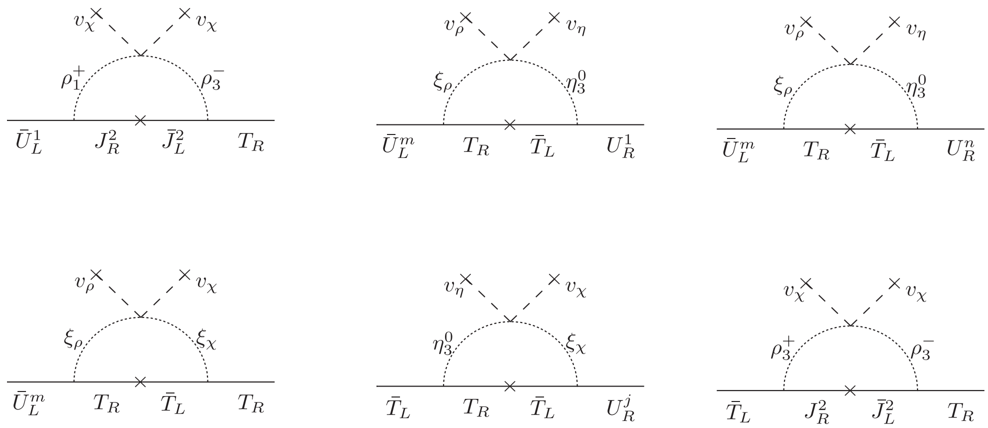

FIG. 1. One-loop Feynman diagrams contributing to the entries of the up-type quark mass matrix.

$$(\delta M^U)_{mn} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1} h^U_{\eta 1n} v_\eta v_\rho}{m_T} C_0\left(\frac{m_{\xi_\rho}}{m_T}, \frac{m_{\eta_3^0}}{m_T}\right),$$
$$(\delta M^U)_{m4} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1} h^T_{\chi 11} v_\rho v_\chi}{m_T} C_0\left(\frac{m_{\xi_\rho}}{m_T}, \frac{m_{\xi_\chi}}{m_T}\right), \qquad (39)$$

$$(\delta M^U)_{4j} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\chi 11} h^U_{\eta 1j} v_\eta v_\chi}{m_T} C_0\left(\frac{m_{\eta_3^0}}{m_T}, \frac{m_{\xi_\chi}}{m_T}\right),$$
$$(\delta M^U)_{44} \simeq -\frac{1}{16\pi^2}\frac{h^J_{\rho 12} h^T_{\rho 21} v_\chi^2}{m_{J_2}} D_0\left(\frac{m_{\rho_3^\pm}}{m_{J_2}}\right), \qquad (40)$$

with $m, n = 2, 3$ and $j = 1, 2, 3$. In the above equations, we use the symbol λ to indicate the quartic coupling terms from the scalar potential in (4), and the following functions have been introduced:

$$C_0(\hat{m}_1, \hat{m}_2) = \frac{1}{(1-\hat{m}_1^2)(1-\hat{m}_2^2)(\hat{m}_1^2-\hat{m}_2^2)} \times \left\{\hat{m}_1^2\hat{m}_2^2 \ln\left(\frac{\hat{m}_1^2}{\hat{m}_2^2}\right) - \hat{m}_1^2 \ln \hat{m}_1^2 + \hat{m}_2^2 \ln \hat{m}_2^2\right\}$$
$$D_0(\hat{m}_1) = \lim_{m_2 \to m_1} C_0(\hat{m}_1, \hat{m}_2) = \frac{-1+\hat{m}_1^2 - \ln \hat{m}_1^2}{(1-\hat{m}_1^2)^2}. \qquad (41)$$

By assuming $m_T \sim m_{J^2} \sim v_\chi$ and $m_{\xi_\rho} \sim m_{\xi_\chi} \sim m_{\eta_3^0} \sim m_{\rho_1^\pm} \sim m_{\rho_3^\pm}$, it follows that the most important one-loop correction for vanishing entries of the tree-level up-type quark mass matrix is $(\delta M^U)_{14}$. On the other hand, the one-loop corrections of the nonvanishing entries of the tree-level up-type quark mass matrix can be neglected when compared to their tree-level values. Therefore, the dominant one-loop level contribution to $Tr(M^U(M^U)^\dagger)$ is roughly $|(\delta M^U)_{14}|^2$. Hence, the mass of the up quark can be estimated as

$$m_u \simeq \frac{1}{16\pi^2}\frac{\lambda |h^J_{\rho 12} h^T_{\rho 21}| v_\chi^2}{m_{J_2}} C_0\left(\frac{m_{\rho_1^+}}{m_{J_2}}, \frac{m_{\rho_3^-}}{m_{J_2}}\right). \qquad (42)$$

Therefore, the smallness of the up quark mass can be explained by the loop suppressed radiative seesaw-type process which involves a heavy exotic quark J^2 as well as virtual charged scalars ρ_1^+ and ρ_3^- whose corresponding Yukawa couplings have to be sufficiently small.

B. Down sector

The mass matrix for the down-type quarks in (23) satisfies the following relation:

$$M^D(M^D)^\dagger = \begin{pmatrix} c & 0_{1\times 2} & X_{1n} \\ 0_{2\times 1} & 0_{2\times 2} & 0_{2\times 1} \\ X_{1n}^\dagger & 0_{1\times 2} & Y_{nm} \end{pmatrix}, \qquad (43)$$

with $n, m = 2, 3$, where

$$c = \left[\sum_{i=1}^{3} |h^D_{\rho 1i}|^2 + \sum_{n=2}^{3} |h^J_{\rho 1n}|^2\right] v_\rho^2,$$
$$X_{1n} = \left[\sum_{i=1}^{3} h^D_{\rho 1i}(h^D_{\chi ni})^\dagger + \sum_{m=2}^{3} h^J_{\rho 1m}(h^J_{\chi nm})^\dagger\right] v_\rho v_\chi, \qquad (44)$$
$$Y_{nm} = \left[\sum_{i=1}^{3} h^D_{\chi ni}(h^D_{\chi mi})^\dagger + \sum_{p=2}^{3} h^J_{\chi np}(h^J_{\chi mp})^\dagger\right] v_\chi^2,$$

from which it follows that

$$\det[M^D(M^D)^\dagger] = 0. \qquad (45)$$

Therefore, the mass matrix texture M^D leads to massless down and strange quarks, which is not phenomenological viable. Besides that, the mass matrix $M^D(M^D)^\dagger$ is partially diagonalized by a rotation matrix V_L^D according to

$$(V_L^D)^\dagger M^D(M^D)^\dagger V_L^D \simeq \begin{pmatrix} m_b^2 & 0 & 0 & 0_{1\times 2} \\ 0 & 0 & 0 & 0_{1\times 2} \\ 0 & 0 & 0 & 0_{1\times 2} \\ 0_{2\times 1} & 0_{2\times 1} & 0_{2\times 1} & Y \end{pmatrix}, \qquad (46)$$

where

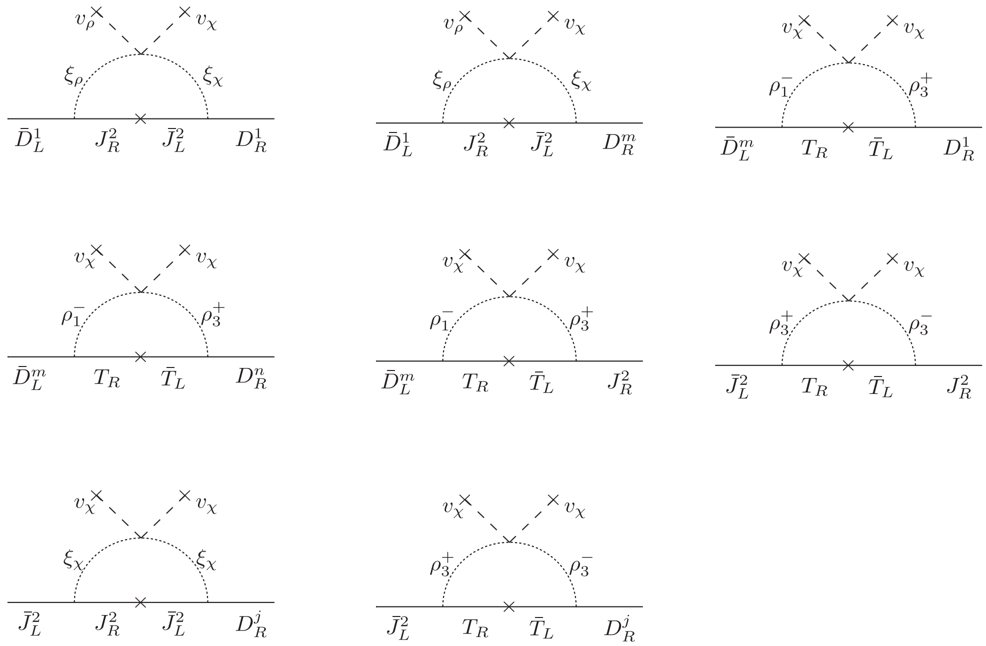

FIG. 2. One-loop Feynman diagrams contributing to the entries of the down-type quark mass matrix.

$$V_L^D = \begin{pmatrix} 1 & 0_{1\times2} & F \\ 0_{2\times1} & 1_{2\times2} & 0_{2\times1} \\ -F^\dagger & 0_{1\times2} & 1_{2\times2} \end{pmatrix} \tag{47}$$

and

$$m_b^2 \simeq c - 2\sum_{n,m=2}^{3} X_{1n} Y_{nm}^{-1} X_{1m}^\dagger, \qquad F_{1m} = \sum_{n=2}^{3} X_{1n} Y_{nm}^{-1}. \tag{48}$$

This shows that radiative corrections at one-loop level have to be introduced in order to generate the masses for the down and strange quarks. For the sake of simplicity, we assume a diagonal base for the exotic quarks J_2 and J_3 and the hierarchy $M_{J_3} \gg M_{J_2}$, which suppress the mixing terms with the J_3 quark at low energy. These assumptions imply that the masses of the exotic quarks J^2 and J^3 are given by

$$M_{J_2} = v_\chi h^J_{\chi 22}, \qquad M_{J_3} = v_\chi h^J_{\chi 33}. \tag{49}$$

Some entries of the down-type quark mass matrix receive loop corrections involving neutral scalars ξ_ρ, ξ_χ with the heavy exotic quark J^2, and charged scalars $\rho_1^\pm$ and $\rho_3^\pm$ with heavy exotic quark T running in the internal lines of the loops. These virtual scalars couple to real neutral scalars due to the scalar quartic interactions. The leading one-loop level contributions to the entries of the down-type quark mass matrix come from the Feynman diagrams shown in Fig. 2. After these radiative corrections are taken into account, the down-type quark mass matrix takes the following form:

$$M^D = \begin{pmatrix} v_\rho h^D_{\rho 11} + (\delta M^D)_{11} & v_\rho h^D_{\rho 12} + (\delta M^D)_{12} & v_\rho h^D_{\rho 13} + (\delta M^D)_{13} & v_\rho h^J_{\rho 12} & 0 \\ (\delta M^D)_{22} & (\delta M^D)_{23} & 0 & (\delta M^D)_{24} & 0 \\ (\delta M^D)_{32} & (\delta M^D)_{33} & 0 & (\delta M^D)_{34} & 0 \\ v_\chi h^D_{\chi 21} + (\delta M^D)_{41} & v_\chi h^D_{\chi 22} + (\delta M^D)_{42} & v_\chi h^D_{\chi 23} + (\delta M^D)_{43} & v_\chi h^J_{\chi 22} + (\delta M^D)_{44} & 0 \\ 0 & 0 & 0 & 0 & v_\chi h^J_{\chi 33} \end{pmatrix}, \tag{50}$$

where their dominant loop corrections are given by

$$(\delta M^D)_{11} \simeq -\frac{1}{16\pi^2} \frac{\lambda h^J_{\rho 12} h^D_{\chi 21} v_\rho v_\chi}{m_{J_2}} C_0\left(\frac{m_{\xi_\rho}}{m_{J_2}}, \frac{m_{\xi_\chi}}{m_{J_2}}\right), \qquad (\delta M^D)_{1m} \simeq -\frac{1}{16\pi^2} \frac{\lambda h^J_{\rho 12} h^D_{\chi 2m} v_\rho v_\chi}{m_{J_2}} C_0\left(\frac{m_{\xi_\rho}}{m_{J_2}}, \frac{m_{\xi_\chi}}{m_{J_2}}\right), \tag{51}$$

$$(\delta M^D)_{m1} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1}h^D_{\rho 11}v_\chi^2}{m_T}C_0\left(\frac{m_{\rho_1^-}}{m_T},\frac{m_{\rho_3^+}}{m_T}\right), \qquad (\delta M^D)_{mn} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1}h^D_{\rho 1n}v_\chi^2}{m_T}C_0\left(\frac{m_{\rho_1^-}}{m_T},\frac{m_{\rho_3^+}}{m_T}\right), \tag{52}$$

$$(\delta M^D)_{m4} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho m1}h^J_{\rho 12}v_\chi^2}{m_T}C_0\left(\frac{m_{\rho_1^-}}{m_T},\frac{m_{\rho_3^+}}{m_T}\right), \qquad (\delta M^D)_{44} \simeq -\frac{1}{16\pi^2}\frac{\lambda h^T_{\rho 21}h^J_{\rho 12}v_\chi^2}{m_T}D_0\left(\frac{m_{\rho_3^\pm}}{m_T}\right), \tag{53}$$

$$(\delta M^D)_{4j} \simeq -\frac{1}{16\pi^2}\left[\frac{\lambda h^J_{\chi 22}h^D_{\chi 2j}v_\chi^2}{m_{J_2}}D_0\left(\frac{m_{\xi_\chi}}{m_{J_2}}\right)+\frac{\lambda h^T_{\rho 21}h^J_{\rho 1j}v_\chi^2}{m_T}D_0\left(\frac{m_{\rho_3^\pm}}{m_T}\right)\right], \tag{54}$$

with $m, n = 2, 3$ and $j = 1, 2, 3$. By assuming $m_T \sim m_{J^2} \sim v_\chi$ and $m_{\xi_\rho} \sim m_{\xi_\chi} \sim m_{\rho_1^\pm} \sim m_{\rho_3^\pm}$, it follows that the two most important loop corrections to $Tr(M^D(M^D)^\dagger)$ come from terms of the order $\sum_{m=2}^3 |(\delta M^D)_{m4}|^2$ and $\sum_{m=2}^3\sum_{n=2}^3 |(\delta M^D)_{mn}|^2$. These terms come from the loop corrections of the tree-level vanishing entries of the down-type quark mass matrix. On the other hand, as in the up sector, the one-loop corrections for the nonvanishing entries of the tree-level down-type quark mass matrix can be neglected when compared to their tree values. Therefore, for the case $|h^D_{\rho 1n}| < <|h^J_{\rho 12}|$ with $n = 1, 2$, the masses of the down and strange quarks can be estimated as

$$m_d \simeq \frac{1}{16\pi^2}\frac{\lambda v_\chi^2}{m_T}\sqrt{\sum_{m=2}^3\sum_{n=2}^3 |h^T_{\rho m1}h^D_{\rho 1n}|^2 C_0\left(\frac{m_{\rho_1^-}}{m_T},\frac{m_{\rho_3^+}}{m_T}\right)}, \qquad m_s \simeq \frac{1}{16\pi^2}\frac{\lambda |h^J_{\rho 12}| v_\chi^2}{m_T}\sqrt{\sum_{m=2}^3 |h^T_{\rho m1}|^2 C_0\left(\frac{m_{\rho_1^-}}{m_T},\frac{m_{\rho_3^+}}{m_T}\right)}. \tag{55}$$

We can see that the charged scalar loop contributions are crucial to give masses to the down and strange quarks. Besides that, the lightness of the down quark can be explained from the smallness of $|h^D_{\rho 1n}|$ as well as from the loop suppressed radiative seesaw-type process which involves a heavy exotic quark T as well as virtual charged scalars ρ_1^- and ρ_3^+. Furthermore, the inequality $|h^D_{\rho 1n}| < <|h^J_{\rho 12}|$ can explain the hierarchy between the down and strange quark masses.

V. CONCLUSIONS

In this paper, we discuss the generation of quark masses in a model based on the gauge symmetry $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ where this symmetry is spontaneously broken to the SM electroweak group $SU(2)_L \otimes U(1)_Y$ at the TeV scale. The Abelian nonuniversal $U(1)_X$ symmetry in the quark sector exhibited in this 3-3-1 model leads to the tree-level cancellation of the Yukawa couplings not allowed by the symmetry. Indeed, since the $U(1)_X$ symmetry of the model distinguishes one family from the other two, the zero-texture structures obtained by (11) and (12) arise naturally, which will lead to only one family (the third) obtaining tree-level masses. The $U(1)_X$ quantum numbers for the exotic quarks T and J are obtained by the condition of cancellation of anomalies, which leads to the mixing terms shown in Eq. (18). These mixing couplings will produce a tree-level mass for the middle quark (charm quark), while the lighter quarks remain massless due to the symmetry. Thus, it is necessary to generate radiative corrections involving scalars and exotic quarks in the internal lines in order to obtain the complete mass spectrum. In this framework, we assume that the CP-odd neutral scalars are much heavier than the heavy exotic quarks T and J^2, and we restrict to the scenario characterized by the absence of mixing between the heavy exotic quark J^3 and the remaining down-type quarks. We find that the mixings between the SM quarks and the exotic quarks as well as the seesaw-type radiative mechanism are crucial to explain the hierarchy of the quark mass spectrum.

ACKNOWLEDGMENTS

R. M. and F. O. thank Colciencias for financial support.

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