Hierarchical fermions and detectable Z' from effective two-Higgs-triplet 3-3-1 model

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(Received 22 January 2018; published 29 March 2018)

We develop a $\mathrm{SU}(3)_C \otimes \mathrm{SU}(3)_L \otimes \mathrm{U}(1)_X$ model where the number of fermion generations is fixed by cancellation of gauge anomalies, being a type of 3-3-1 model with new charged leptons. Similarly to the economical 3-3-1 models, symmetry breaking is achieved effectively with two scalar triplets so that the spectrum of scalar particles at the TeV scale contains just two *CP* even scalars, one of which is the recently discovered Higgs boson, plus a charged scalar. Such a scalar sector is simpler than the one in the Two Higgs Doublet Model, hence more attractive for phenomenological studies, and has no flavor changing neutral currents (FCNC) mediated by scalars except for the ones induced by the mixing of Standard Model (SM) fermions with heavy fermions. We identify a global residual symmetry of the model which guarantees mass degeneracies and some massless fermions whose masses need to be generated by the introduction of effective operators. The fermion masses so generated require less fine-tuning for most of the SM fermions and FCNC are naturally suppressed by the small mixing between the third family of quarks and the rest. The effective setting is justified by an ultraviolet completion of the model from which the effective operators emerge naturally. A detailed particle mass spectrum is presented, and an analysis of the Z' production at the LHC run II is performed to show that it could be easily detected by considering the invariant mass and transverse momentum distributions in the dimuon channel.

DOI: 10.1103/PhysRevD.97.055047

I. INTRODUCTION

The measurements of the Higgs boson properties and their actual agreement with the Standard Model predictions [1–4] have corroborated the simplest implementation of the Higgs mechanism as the source of electroweak symmetry breaking. Although the CERN Large Hadron Collider (LHC) has not yet provided clear evidence for new physics, the Standard Model (SM) consolidation has helped us to put in a firmer footing a series of its theoretical problems such as the severe hierarchy of the Yukawa couplings (the flavor problem), including the neutrino masses and mixing problem; the number of fermion generations; the chiral nature of the electroweak interaction; matter-antimatter asymmetry of the Universe; the strong *CP* problem; the dark matter content of the Universe; and the vacuum stability. The seeking of solutions for one or more of these problems has often guided the development of new models by extending the field content of the SM or, sometimes simultaneously, enlarging its symmetries.

Concerning the empirical observation of just three generations of fermions, 3-3-1 models offer a plausible explanation [5–11]. In these theoretical constructions the $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ symmetry group of the electroweak interactions is extended to $\mathrm{SU}(3)_L \otimes \mathrm{U}(1)_X$, in such a way that cancellation of all gauge anomalies involves necessarily all the three fermion generations. As it happens, there are different types of 3-3-1 models depending on the matter content fixed by a parameter β in the electric charge operator

$$Q = T_3 + \beta T_8 + XI, \tag{1}$$

where T_3 and T_8 are the diagonal generators of $\mathrm{SU}(3)_L$ built as $T_a = \frac{\lambda_a}{2}$ from the Gell-Mann matrices λ_a, with $a = 1, \ldots, 8$; and X refers to the $\mathrm{U}(1)_X$ charge. Standard Model left-handed lepton fields take part in $\mathrm{SU}(3)_L$ triplets, $\psi_{iL} = (\nu_i e_i^- E_i^{q_E})_L^T$, having $X_\psi = -\frac{1}{2}(1 + \frac{\beta}{\sqrt{3}})$. The third

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components, $E_i^{q_E}$, are new lepton fields with electric charge $q_E = -\frac{1}{2}(1+\sqrt{3}\beta)$. The particular choice $\beta = -\frac{1}{\sqrt{3}}$ leads to models where the new leptons E_{iL} do not carry an electric charge [5,9,11,12]. Other well developed constructions have $\beta = -\sqrt{3}$ and, in this case, E_{iL} carry one unit of electric charge so that it could be identified with the charged antileptons, i.e., $E_{iL} \equiv (l_{iR}^-)^c$ [6,7], or even represent new charged leptons [10]. For other possibilities, see [13].

For all models, the cancellation of gauge anomalies requires two antitriplets (triplets) and one triplet (antitriplet) of left-handed quarks when taking into account triplets (antitriplets) of leptons. This feature leads to flavor changing neutral currents (FCNC) mainly through a vector boson Z' whose mass is related to the energy scale in which the $SU(3)_L \otimes U(1)_X$ symmetry is broken down spontaneously to $SU(2)_L \otimes U(1)_Y$. Bounds on the Z' boson mass have been obtained from the LHC data for the versions with $\beta = \pm 1/\sqrt{3}$, $-\sqrt{3}$ in Ref. [14], and recent analyses on FCNC have been performed in Refs. [15–17].

Many works have been published exploring the theoretical and phenomenological benefits of these models, showing that they are good candidates for describing new physics. For example, it is possible to include supersymmetry in such a context [18–21], as well as to construct left-right extensions [22], which have recently been subject of some studies [23–26].

To implement spontaneous symmetry breaking, three or more scalar multiplets getting vacuum expectation value (vev) at the GeV-TeV scale have been considered in 3-3-1 models. As a consequence, the scalar potential has many free parameters, being more complex than, e.g., the Two Higgs Doublet Model (see Ref. [27] for a review). However, it has been shown that it is possible to break the symmetries down in some 3-3-1 models by taking into account two scalar triplets only [17,28–31]. These constructions are phenomenologically attractive once they have a simpler scalar potential, predicting only three Higgs bosons. With the introduction of effective operators, masses for all fermions can be generated [17,28].

Our aim in this work is to develop a version of the 3-3-1 model with $\beta = \frac{1}{\sqrt{3}}$ distinct from its first proposals [32,33] and from other similar models focused on neutrino masses and mixing [34,35]. We focus on this model because the versions with $\beta = \pm\sqrt{3}$ [6,7,10] become strongly interacting at an energy of few TeV [7,36–39]. We comment on the case $\beta = -1/\sqrt{3}$ when relevant. We show that a consistent symmetry breaking pattern is obtained for this model with only two scalar triplets getting vev. Also, we present a simple mechanism where the low energy effective operators required to generate mass for some fermions arise after the integration of a supposedly heavy scalar triplet. These effective operators share similarities with those in the Froggatt-Nielsen mechanism in the sense that they generate more natural, less fine-tuned, masses for most of the fermions of the model, when compared to the mass generation in the Standard Model. The scalar particle spectrum of the model is composed of just two neutral CP even scalars, with one of them directly identified with the discovered 125 GeV Higgs boson, plus a charged one and its antiparticle. Also, five new vector bosons, V^+, V^-, V^0, $V^{0\dagger}$, and Z', are predicted by the model. A study of the production signals at the LHC of the Z' boson is performed.

The paper is organized as follows: in Sec. II, the essential aspects of the model are presented, including symmetry breakdown, residual symmetries, and the particle spectra of scalar and vector bosons; Fermion masses are treated in Sec. III; flavor changing interactions are analyzed in Sec. IV; in Sec. V, a simple UV-completion able to generate the needed effective operators for fermion masses is discussed; the Z' boson phenomenology is presented in Sec. VI; and our conclusions are given in Sec. VII.

II. THE MODEL

We focus on the model with $\beta = 1/\sqrt{3}$ in Eq. (1). Therefore, the left-handed lepton fields form $SU(3)_L$ triplets, with the right-handed lepton fields in $SU(3)_L$ singlets, as follows:

$$\psi_{iL} = (\nu_i, e_i^-, E_i^-)_L^T \sim (\mathbf{1}, \mathbf{3}, -2/3),$$
$$\nu_{iR} \sim (\mathbf{1}, \mathbf{1}, 0), \qquad e'^-_{sR} \sim (\mathbf{1}, \mathbf{1}, -1), \tag{2}$$

where $i = 1, 2, 3$, is the generation index, and $s = 1, \ldots, 6$, with $e'^-_{sR} \equiv (e^-_{iR}, E^-_{iR})$. The numbers in parentheses refer to the field transformation properties under $SU(3)_C$, $SU(3)_L$, and $U(1)_X$, respectively. We consider the right-handed neutrino fields, ν_{iR}, in order to generate small masses to the left-handed neutrinos through the usual seesaw mechanism. The fields E^-_{iL}, required to complete the $SU(3)_L$ representation, along with the right-handed components, give rise to three *heavy leptons*.

Given the above lepton multiplets, as first observed long ago, gauge anomalies are canceled when the three families of quarks are included nonuniversally into two antitriplets and one triplet of $SU(3)_L$ for the left-handed parts, and the corresponding right-handed fields assigned to singlets:

$$Q_{aL} = (d_a, -u_a, U_a)_L^T \sim (\mathbf{3}, \mathbf{3}^*, 1/3),$$
$$Q_{3L} = (u_3, d_3, D)_L^T \sim (\mathbf{3}, \mathbf{3}, 0),$$
$$u'_{mR} \sim (\mathbf{3}, \mathbf{1}, 2/3), \qquad d'_{nR} \sim (\mathbf{3}, \mathbf{1}, -1/3), \tag{3}$$

where $a = 1,\ 2,\ m = 1, \ldots, 5,\ n = 1, \ldots, 4$, with $u'_{mR} \equiv (u_{iR}, U_{aR})$ and $d'_{nR} \equiv (d_{iR}, D_R)$. Besides the quark fields of the SM, this model has two extra up-type quark fields, U_a, and one down-type field, D. Such fields, as well as E_i, get their masses at the energy scale w, in which the $SU(3)_L \otimes U(1)_X$ is supposedly broken down to $SU(2)_L \otimes U(1)_Y$. Once that energy scale must be higher than the electroweak scale, i.e., $w > v = 246$ GeV, it is natural for the new elementary fermions associated with those fields to be heavier than the standard ones.

As we have already mentioned, the set of fields in Eqs. (2) and (3) is such that the cancellation of gauge anomalies involves the three fermion generations. This contrasts with the SM where the cancellation of anomalies occurs in each family, independently.

In principle, the choice of which generation of left-handed quark is assigned to a triplet is arbitrary. But, the fact that not all left-handed quark multiplets have the same transformation properties leads to new sources of FCNC. This has been explored in various works considering different versions of 3-3-1 models. Constructions with the third generation transforming differently from the first two are less restricted by bounds of processes involving FCNC. We show in Sec. IV that FCNC interactions are naturally suppressed in our model due to its peculiar mass generation mechanism for the fermions.

The following two scalar triplets realize the spontaneous breaking of the $\mathrm{SU}(3)_L \otimes \mathrm{U}(1)_X$ symmetry down to $\mathrm{U}(1)_Q$ of the electromagnetic interactions:

$$\rho \equiv (\rho_1^0 \rho_2^- \rho_3^-)^T \sim (\mathbf{1}, \mathbf{3}, -2/3),$$
$$\chi \equiv (\chi_1^+ \chi_2^0 \chi_3^0)^T \sim (\mathbf{1}, \mathbf{3}, 1/3). \tag{4}$$

This is the minimal set of scalar fields that can perform the required symmetry breakdown.

From the fermionic and the scalar multiplets in Eqs. (2), (3), and (4), we write down the following Yukawa Lagrangian:

$$\begin{aligned} -\mathcal{L}_Y = {} & h^E_{is}\overline{\psi_{iL}}\chi e'_{sR} + h^\nu_{ij}\overline{\psi_{iL}}\rho\nu_{jR} + \frac{1}{2} m_{ij}\overline{(\nu_{iR})^c}\nu_{jR} \\ & + h^U_{am}\overline{Q_{aL}}\chi^* u'_{mR} + h^d_{an}\overline{Q_{aL}}\rho^* d'_{nR} + f^u_m\overline{Q_{3L}}\rho u'_{mR} \\ & + f^D_n\overline{Q_{3L}}\chi d'_{nR} + \text{H.c.}, \end{aligned} \tag{5}$$

where the complex coupling constants are such that: h^E_{is} is a 3×6 matrix; h^ν_{ij}, is a 3×3 matrix; $m_{ij} = m_i\delta_{ij}$ is a 3×3 diagonal matrix; h^U_{am} and h^d_{an} are 2×5 and 2×4 matrices, respectively; f^u_m and f^D_n are 1×5 and 1×4 matrices, respectively.

With only the two scalar triplets in Eq. (4), the most general renormalizable scalar potential is simply given by

$$\begin{aligned} V(\chi,\rho) = {} & \mu_1^2\rho^\dagger\rho + \mu_2^2\chi^\dagger\chi + \lambda_1(\rho^\dagger\rho)^2 + \lambda_2(\chi^\dagger\chi)^2 \\ & + \lambda_3(\chi^\dagger\chi)(\rho^\dagger\rho) + \lambda_4(\chi^\dagger\rho)(\rho^\dagger\chi). \end{aligned} \tag{6}$$

We assume that the quadratic mass parameters, $\mu_{1,2}^2 < 0$, and the self-interaction coupling constants, λ_i, $i = 1, \ldots 4$, are such that the scalar fields will develop nonvanishing vevs, $\langle\chi\rangle$, $\langle\rho\rangle \neq 0$.

TABLE I. Charges of the Abelian symmetries of the model.

#	χ	ρ	ψ_{iL}	(e_{iR}, E_{iR})	ν_{iR}	Q_{aL}	Q_{3L}	(d_{iR}, D_{3R})	(u_{iR}, U_{aR})
X	1/3	−2/3	−2/3	−1	0	1/3	0	−1/3	2/3
PQ	1	1	1	0	0	−1	1	0	0
B	0	0	0	0	0	1/3	1/3	1/3	1/3
Lep	0	0	1	1	1	0	0	0	0

A. Symmetry breaking and residual symmetries

Besides being invariant under the gauge symmetries $\mathrm{SU}(3)_C\otimes\mathrm{SU}(3)_L\otimes\mathrm{U}(1)_X$, our model presents invariance under certain global symmetries that we will now describe. When only renormalizable operators are taken into account, and right-handed neutrinos are not introduced, one can check that the Lagrangian is invariant under three extra global U(1) symmetries. Two of which can be taken as the Baryon and the Lepton number symmetries ($\mathrm{U}(1)_B$ and $\mathrm{U}(1)_{\mathrm{Lep}}$), while the other one, being associated with a $[\mathrm{SU}(3)_C]^2\times\mathrm{U}(1)$ anomaly, is a Peccei-Quinn-like symmetry ($\mathrm{U}(1)_{PQ}$).[1] The latter symmetry is the continuous version of the center of $\mathrm{SU}(3)_L$ in which every triplet, **3**, carries unit charge while every antitriplet, $\mathbf{3}^*$, carries the opposite charge. This is possible in our case due to the absence of any trilinear couplings of the form $\mathbf{3}\times\mathbf{3}\times\mathbf{3}$. Therefore, when the scalar fields get vev, the $\mathrm{U}(1)_{PQ}$ symmetry will be broken, but its charges will be part of another remaining global symmetry as we show below. In Table I the quantum numbers associated with the U(1) symmetries are presented for all the matter fields.

Thus, our model is actually invariant under a larger symmetry group: $\mathrm{SU}(3)_C\otimes\mathrm{SU}(3)_L\otimes\mathrm{U}(1)_X\otimes\mathrm{U}(1)_{PQ}\otimes\mathrm{U}(1)_B(\otimes\mathrm{U}(1)_{\mathrm{Lep}})$ with the additional Abelian symmetries being global ones. We put $\mathrm{U}(1)_{\mathrm{Lep}}$ in parentheses to call the reader's attention to the fact that such a symmetry is only present when the bare Majorana mass term for the ν_{iR} singlets in Eq. (5) is absent.

Since scalar fields transform trivially under the $\mathrm{SU}(3)_C\otimes\mathrm{U}(1)_B(\otimes\mathrm{U}(1)_{\mathrm{Lep}})$ subgroup, there is no way in which nonvanishing vevs of the scalar fields will trigger spontaneous symmetry breaking of such a structure. Consequently, we neglect this subgroup for now on and focus only on the groups affected by spontaneous symmetry breaking, i.e. $\mathrm{SU}(3)_L\otimes\mathrm{U}(1)_X\otimes\mathrm{U}(1)_{PQ}$, containing ten independent generators.

We begin our analysis by noting that the electric charge operator in Eq. (1) with $\beta = 1/\sqrt{3}$ implies that the second

[1]We mean here that such a global symmetry is chiral and anomalous, but it will be broken explicitly, as discussed in Sec. III, implying that the Peccei-Quinn mechanism does not take place in our model. Implementations of the Peccei-Quinn symmetry for the strong CP problem in the 3-3-1 models can be found, for example, in Refs. [40,41].

and third component fields of any $SU(3)_L$ triplet or antitriplet always have the same electric charge. This means that the charge operator is invariant by

$$U^\dagger Q U = Q, \tag{7}$$

where U can be an arbitrary SU(2) transformation in the 2-3 sub-block or any diagonal transformation.[2] It is the former group which is special to this case, and we denote this group as $SU(2)_{rep}$ where rep stands for reparametrization. The situation is different from many gauge extensions of the SM in which there are no equal charge fields in the same multiplet, and the only transformations that leave Q invariant are the diagonal ones which include Q itself, modulo Abelian factors. This is also different from horizontal spaces that are present even before symmetry breaking, such as in the Two Higgs Doublet extension of the SM or the extension (UV completion) of the current model with one more scalar triplet; see Sec. V. Reparametrization symmetry means that we can rotate all the fields of the theory by an $SU(2)_{rep}$ transformation, including the vev in Eq. (8), without affecting the physical content. The physical invariance is ensured because $SU(2)_{rep}$ is a subgroup of the original $SU(3)_L$ global gauge group.

Therefore, without loss of generality we can consider that the minimum of the potential in Eq. (6) is attained at the vevs

$$\langle\chi\rangle = \frac{1}{\sqrt{2}}(0,0,w)^T, \tag{8}$$

$$\langle\rho\rangle = \frac{1}{\sqrt{2}}(v,0,0)^T. \tag{9}$$

If we had considered the more general vev $\langle\chi\rangle = (0,v',w')^T/\sqrt{2}$, the reparametrization symmetry in (7) would allow us to rotate $\langle\chi\rangle$ to the original form in Eq. (8) without affecting $\langle\rho\rangle$ in Eq. (9). Consequently, the fields are also transformed so that Eqs. (8) and (9) can be taken from the start. A direct consequence is that the vector bosons $V^\pm$ and $W^\pm$ do not mix at tree level. It has to be pointed out that the reparametrization symmetry applies to other models, like the 3-3-1 model defined by $\beta = -1/\sqrt{3}$, having a scalar triplet with two neutral components that could acquire vevs, so that we can make a rotation in order to have just one component with vev.

Given that we are considering $w > v$, the spontaneous symmetry breaking induced by the vevs in Eqs. (8) and (9) happens in two stages: first with $\langle\chi\rangle$ realizing the breakdown

$$\begin{aligned} &SU(3)_L \otimes U(1)_X \otimes U(1)_{PQ} \\ &\quad \to SU(2)_L \otimes U(1)_Y \otimes U(1)_{PQ'}, \end{aligned} \tag{10}$$

with the hypercharge given by $Y = T_8/\sqrt{3} + X$, and the charges of the global symmetry $PQ' = 3X - PQ$; and second with $\langle\rho\rangle$ realizing the breakdown

$$SU(2)_L \otimes U(1)_Y \otimes U(1)_{PQ'} \to U(1)_Q \otimes U(1)_{\mathcal{G}}, \tag{11}$$

with Q the electric charge operator of Eq. (1) and charges of a global symmetry given by[3]

$$\mathcal{G} = 2T_3 + X - \frac{1}{3}PQ. \tag{12}$$

It is easy to see that this operator is unbroken by the vevs of Eqs. (8) and (9) when we write it explicitly for χ and ρ: $\mathcal{G}(\chi) = \text{diag}(1,-1,0)$ and $\mathcal{G}(\rho) = \text{diag}(0,-2,-1)$.[4] Thus, two independent generators out of the ten initial remain unbroken after spontaneous symmetry breaking. The eight would-be Goldstone bosons associated with the broken generators are all absorbed to form the longitudinal degrees of freedom of the massive vector bosons: Z, $W^\pm$, Z', $V^\pm$, and the neutral non-Hermitian V^0 and $V^{0\dagger}$, with $\mathcal{G}$ charges 0, ±2, 0, ±1, -1 and $+1$, respectively; see Eq. (30). Consequently, from the twelve degrees of freedom contained in the two scalar triplets, four are left as physical scalar boson fields: two neutral CP even, h and H, plus the charged ones $\varphi^\pm$. The mass spectra for the scalar bosons and vector bosons are shown in the next section. We anticipate some mass degeneracy from the conservation of $\mathcal{G}$. For example, since V^0 and $V^{0\dagger}$ are the only neutral gauge bosons with $\mathcal{G}$ charges ∓ 1, we expect that they remain mass degenerate and do not split into two neutral gauge bosons with different masses. This expectation is confirmed by the explicit calculation of the mass matrices; see Sec. II C.

The $U(1)_{\mathcal{G}}$ symmetry also has the property of being chiral for the second components of fermion triplets (antitriplets) and their right-handed counterparts in singlets of $SU(3)_L$. As a result the standard charged leptons, two up-type quarks and one down-type quark, are left massless even at the perturbative level, as pointed out in [42]. In order to overcome this problem we introduce in Sec. III dimension-5 operators which explicitly break the $U(1)_{PQ}$ and, consequently, the $U(1)_{\mathcal{G}}$ symmetry. As we will see, such operators can be generated at low energies in an UV-completed model with a heavy scalar triplet which is integrated out. In such a setting the $\mathcal{G}$ charge is only

[2]Obviously, we can perform a more general reparametrization in $SU(3)_L$ without physical consequence but the charge operator Q would change.

[3]If both neutral components of χ acquire a vev, i.e. $\sqrt{2}\langle\chi\rangle = (0,u,w)$, there still remains a conserved symmetry generated by $\mathcal{G}_\theta = 2(1-2\sin^2\theta)T_3 + \sin(2\theta)T_6 + (1-3\sin^2\theta)X - \frac{1}{3}PQ$, with $\tan\theta = u/w$ for the case that u, w are real and positive. If they were complex a more general expression can be written reparametrized by $SU(2)_{rep}$.

[4]It is also clear that $U(1)_{\mathcal{G}}$ would be generally broken if there is an additional scalar triplet that can acquire a vev in its second component.

broken by the soft breaking of PQ, and thus it remains approximately conserved.

B. Scalar bosons

In order to find the scalar field masses and corresponding physical states, let us first write the scalar triplets as

$$\chi = \begin{pmatrix} \chi_1^+ \\ \chi_2^0 \\ \frac{1}{\sqrt{2}}(w + S_3 + iA_3) \end{pmatrix} \quad \text{and}$$

$$\rho = \begin{pmatrix} \frac{1}{\sqrt{2}}(v + S_1 + iA_1) \\ \rho_2^- \\ \rho_3^- \end{pmatrix}, \tag{13}$$

where we have decomposed the neutral fields which acquire a nonvanishing vev into scalar and pseudoscalar contributions, S_i and A_i, respectively. In the approximation that the global charge $\mathcal{G}$ is exactly conserved, we can expect from its conservation that both χ_2^0 and ρ_2^- already have definite masses (they are would-be Goldstone bosons) and the possible pairs that can mix are (χ_1^+, ρ_3^+), (S_1, S_3) and (A_1, A_3), assuming CP conservation.

The minimum condition for the potential leads to the constraint equations

$$\mu_1^2 + \lambda_1 v^2 + \frac{1}{2}\lambda_3 w^2 = 0$$
$$\mu_2^2 + \lambda_2 w^2 + \frac{1}{2}\lambda_3 v^2 = 0, \tag{14}$$

from which the quadratic mass parameters $\mu_{1,2}^2$ can be eliminated.

The mass matrix, derived from the scalar potential, for the CP even scalars in the basis (S_1, S_3) is

$$M_0^2 = \begin{pmatrix} 2\lambda_1 v^2 & \lambda_3 vw \\ \lambda_3 vw & 2\lambda_2 w^2 \end{pmatrix}. \tag{15}$$

This leads to the quadratic mass eigenvalues

$$m_h^2 = \lambda_1 v^2 + \lambda_2 w^2 - \sqrt{\lambda_3^2 v^2 w^2 + (\lambda_2 w^2 - \lambda_1 v^2)^2}, \tag{16}$$

$$m_H^2 = \lambda_1 v^2 + \lambda_2 w^2 + \sqrt{\lambda_3^2 v^2 w^2 + (\lambda_2 w^2 - \lambda_1 v^2)^2}, \tag{17}$$

corresponding, respectively, to the mass eigenstates

$$h = \cos\theta S_1 + \sin\theta S_3 \tag{18}$$

$$H = -\sin\theta S_1 + \cos\theta S_3, \tag{19}$$

where $\tan 2\theta = \lambda_3 vw/(\lambda_2 w^2 - \lambda_1 v^2)$. We identify h as the state corresponding to the observed Higgs boson with mass of 125 GeV. In the limit $\theta \to 0$, the tree level couplings of h to the electroweak vector bosons W and Z are the same as the Standard Model Higgs boson.

The particle spectrum of the model does not contain CP odd neutral scalar fields. The pseudoscalar fields A_1 and A_3 are absorbed in the massive vector bosons Z and Z'. In particular, the complex field $\chi_2^0 = (S_2 + iA_2)/\sqrt{2}$ in the triplet χ does not get a mass term and plays the role of the Goldstone boson absorbed in the non-Hermitian neutral vector boson V^0 (both have $\mathcal{G}$ charge -1). This contrasts with the Two Higgs Doublet Models, which necessarily contain a neutral pseudoscalar in the particle spectrum.

For the charged scalar fields, it can be seen that only ρ_3^+ and χ_1^+ mix with each other so that their mass matrix, in the basis $(\rho_3^\pm, \chi_1^\pm)$, is

$$M_\pm^2 = \frac{\lambda_4}{2}\begin{pmatrix} w^2 & vw \\ vw & v^2 \end{pmatrix}, \tag{20}$$

whose the nonzero eigenvalue

$$m_{\varphi^\pm}^2 = \frac{\lambda_4}{2}(v^2 + w^2), \tag{21}$$

corresponds to the squared mass of a charged scalar state given by

$$\varphi^\pm = \frac{1}{\sqrt{v^2 + w^2}}(w\rho_3^\pm + v\chi_1^\pm). \tag{22}$$

The orthogonal eigenstates $G_{31}^\pm = (v\rho_3^\pm - w\chi_1^\pm)/\sqrt{v^2 + w^2}$ and $G_2^\pm = \rho_2^\pm$ are Goldstone bosons which are absorbed to form the longitudinal components of the vector bosons $V^\pm$ and $W^\pm$.

Thus, we see that four, from the initial twelve, degrees of freedom contained in the two scalar triplets remain as the physical scalars h, H, $\varphi^\pm$. The other eight degrees of freedom become the Goldstone modes needed to give mass to the vector bosons $W^\pm$, Z, $V^\pm$, V^0, $V^{0\dagger}$, and Z'. So ρ contains predominantly the SM Higgs h within the $\mathrm{SU}(2)_L$ doublet and the heavy charged Higgs φ^- in its third component while χ contains predominantly the heavy Higgs H in the third component and a small admixture of the φ^+ within the $\mathrm{SU}(2)_L$.

C. Vector bosons

As usual, to determine the physical gauge bosons and their masses, we look at the covariant derivative terms for the scalar fields:

$$\mathcal{L} \supset (D_\mu\rho)^\dagger(D^\mu\rho) + (D_\mu\chi)^\dagger(D^\mu\chi), \tag{23}$$

in which the covariant derivative is defined as

$$D_\mu = \partial_\mu - igW^a_\mu T^a - ig_X X B_\mu = \partial_\mu - iP_\mu, \quad (24)$$

where T^a, with $a = 1, \ldots, 8$, are the $SU(3)_L$ generators as defined in Eq. (1), and X denotes the $U(1)_X$ charge of the field on which D_μ acts; g, g_X are the gauge coupling constants related to $SU(3)_L$ and $U(1)_X$, respectively. The gauge coupling constant g is the same as in the Standard Model, since in 3-3-1 models the gauge group $SU(2)_L$ is totally embedded in $SU(3)_L$. Additionally, the gauge coupling constants are related to the Standard Model electroweak mixing angle θ_W according to

$$t^2 = \frac{g_X^2}{g^2} = \frac{\sin^2\theta_W}{1 - \frac{4}{3}\sin^2\theta_W}. \quad (25)$$

Then, the P_μ matrix for **3** can be written as

$$P_\mu = \frac{g}{2}\begin{pmatrix} W_{3\mu} + \frac{W_{8\mu}}{\sqrt{3}} + 2tB_\mu X & \sqrt{2}W^+_\mu & \sqrt{2}V^+_\mu \\ \sqrt{2}W^-_\mu & -W_{3\mu} + \frac{W_{8\mu}}{\sqrt{3}} + 2tB_\mu X & \sqrt{2}V^0_\mu \\ \sqrt{2}V^-_\mu & \sqrt{2}V^{0\dagger}_\mu & -\frac{2W_{8\mu}}{\sqrt{3}} + 2tB_\mu X \end{pmatrix}, \quad (26)$$

where we have defined the following fields

$$W^+_\mu = \frac{W_{1\mu} - iW_{2\mu}}{\sqrt{2}}, \quad (27)$$

$$V^+_\mu = \frac{W_{4\mu} - iW_{5\mu}}{\sqrt{2}}, \quad (28)$$

$$V^0_\mu = \frac{W_{6\mu} - iW_{7\mu}}{\sqrt{2}}. \quad (29)$$

The $\mathcal{G}$ charge carried by these gauge fields is given by $2T_3$ and yields

$$\mathcal{G}(P_\mu) = \begin{pmatrix} 0 & 2 & 1 \\ & 0 & -1 \\ & & 0 \end{pmatrix}. \quad (30)$$

The vector boson masses arise when the scalar fields in Eq. (23) acquire vevs as in Eqs. (8) and (9). The vector boson fields $W^\pm$, $V^\pm$, V^0, and $V^{0\dagger}$ get the following squared masses

$$M^2_{W^\pm} = \frac{g^2 v^2}{4}, \quad (31)$$

$$M^2_{V^\pm} = \frac{g^2}{4}(v^2 + w^2), \quad (32)$$

$$M^2_{V^0} = M^2_{V^{0\dagger}} = \frac{g^2}{4}w^2. \quad (33)$$

A direct consequence of breaking down the symmetries with just two scalar triplets is the tree level mass splitting prediction $M^2_{V^\pm} - M^2_{V^0} = M^2_W$. Another peculiarity of the model is that a novel sort of neutral current might occur involving V^0, $V^{0\dagger}$, since these vector bosons intermediate transitions between standard and new leptons with the same electric charge.

The gauge boson fields W_3, W_8, and B of the symmetry basis mix with each other leading to the mass matrix, in the basis (W_3, W_8, B),

$$M_0^2 = \frac{g^2}{2}\begin{pmatrix} \frac{v^2}{2} & \frac{v^2}{2\sqrt{3}} & -\frac{2v^2}{3}t \\ \frac{v^2}{2\sqrt{3}} & \frac{(v^2+4w^2)}{6} & \frac{-2(v^2+w^2)}{3\sqrt{3}}t \\ -\frac{2v^2}{3}t & \frac{-2(v^2+w^2)}{3\sqrt{3}}t & \frac{2(4v^2+w^2)}{9}t^2 \end{pmatrix}. \quad (34)$$

The mass eigenstates from this matrix give the photon field, A_μ, and two massive fields, $Z_{1\mu}$ and $Z_{2\mu}$,

$$A_\mu = \frac{\sqrt{3}}{\sqrt{3+4t^2}}\left(tW^3_\mu + \frac{t}{\sqrt{3}}W^8_\mu + B_\mu\right), \quad (35)$$

$$Z_{1\mu} = N_{Z_2}(-3M^2_{Z_2}W^3_\mu + \sqrt{3}(3M^2_{Z_2} - g^2w^2)W^8_\mu + g^2w^2tB_\mu), \quad (36)$$

$$Z_{2\mu} = N_{Z_1}(-3M^2_{Z_1}W^3_\mu + \sqrt{3}(3M^2_{Z_1} - g^2w^2)W^8_\mu + g^2w^2tB_\mu), \quad (37)$$

where the normalization constants are

$$N_{Z_2,Z_1} = [(g^2w^2t)^2 + (3M^2_{Z_2,Z_1})^2 + 3(3M^2_{Z_2,Z_1} - g^2w^2)^2]^{-1/2}.$$

The masses of the neutral vector bosons, $Z_{1\mu}$ and $Z_{2\mu}$, can be written as

$$M^2_{Z_2,Z_1} = \frac{g^2}{18}\Big[(3+4t^2)v^2 + (3+t^2)w^2 \pm \sqrt{-9(3+4t^2)v^2w^2 + ((3+4t^2)v^2 + (3+t^2)w^2)^2}\Big], \quad (38)$$

in such a way that the dominant contributions are

$$M^2_{Z_1} = \frac{g^2v^2}{4\cos^2\theta_W} + \mathcal{O}\left(\frac{v^2}{w^2}\right),$$
$$M^2_{Z_2} = \frac{g^2\cos^2\theta_W w^2}{3-4\sin^2\theta_W} + \mathcal{O}\left(\frac{v^2}{w^2}\right). \quad (39)$$

And there is also the prediction, at tree level, that

$$M^2_{Z_2}/M^2_{V^0} \approx 4\cos^2\theta_W/(3-4\sin^2\theta_W) \approx 1.48, \quad (40)$$

where we have used $\sin^2\theta_W \approx 0.231$.

It is convenient for the discussion on the FCNC in the model to express the mass eigenstates $Z_{2\mu}$ and $Z_{1\mu}$ as linear combinations of the fields Z'_μ and Z_μ, which result from the sequential symmetry breakdown $SU(3)_L \otimes U(1)_X$ and $SU(2)_L \otimes U(1)_Y$, respectively,

$$\begin{pmatrix} Z_1 \\ Z_2 \end{pmatrix} = \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix}\begin{pmatrix} Z \\ Z' \end{pmatrix}. \quad (41)$$

The unitary matrix in Eq. (41) diagonalizes the Z_μ–Z'_μ mass matrix

$$\begin{pmatrix} M^2_Z & M^2_{ZZ'} \\ M^2_{ZZ'} & M^2_{Z'} \end{pmatrix}, \quad (42)$$

where

$$M^2_Z = \frac{g^2v^2}{4\cos^2\theta_W}, \qquad M^2_{ZZ'} = -\frac{M^2_Z}{\sqrt{3-4\sin^2\theta_W}},$$
$$M^2_{Z'} = \frac{M^2_Z + g^2w^2\cos^2\theta_W}{3-4\sin^2\theta_W}. \quad (43)$$

In terms of these elements, the angle θ is

$$\tan(2\theta) = \frac{2M^2_{ZZ'}}{M^2_{Z'} - M^2_Z}. \quad (44)$$

For example, taking $M_{Z'} \approx 4$ TeV, the angle is then $\theta \approx M^2_{ZZ'}/M^2_{Z'} \approx -4\times 10^{-4}$ so that $Z_{2\mu} \approx Z'_\mu$.

Due to the fact that $M^2_{Z_1}$ also depends on the energy scale w related to the breakdown of $SU(3)_L \otimes U(1)_X$ to $SU(2)_L \otimes U(1)_Y$, the model presents a deviation from the Standard Model ρ_0 parameter prediction $\rho_0 = M^2_W/\cos^2\theta_W M^2_{Z^0} = 1$, with M_{Z^0} standing for the Standard Model Z^0 boson mass at tree level. Such a deviation, up to order $(v/w)^2 \ll 1$ in $M^2_{Z_1}$, is given by

$$\Delta\rho_0 \equiv \frac{M^2_W}{\cos^2\theta_W M^2_{Z_1}} - 1 \approx \frac{(v/w)^2}{4\cos^4\theta_W}, \quad (45)$$

where $M^2_{Z_1} = M^2_{Z^0} + \delta M^2_Z$. The actual experimental data furnishes $\Delta\rho_0 \equiv \rho_0 - 1 \lesssim 0.0006$ [43]. Thus, if the tree level contribution is dominant over the radiative corrections we obtain the lower-bound $w \geq 6.5$ TeV, taking into account the value $v = 246$ GeV. For definiteness, we take $w = 10$ TeV which furnishes $M_{Z_2} \approx 4$ TeV, and $M_{V^\pm} \approx M_{V^0} \approx 3.2$ TeV, although in our phenomenological analysis in Sec. VI a lower value for the Z' boson mass is also considered.

III. FERMION MASSES

With the vevs $\langle\chi\rangle$ and $\langle\rho\rangle$ the Yukawa Lagrangian in Eq. (5) leads to the mass matrix for neutrinos, in the basis $(\nu_{iL}, {\nu_{iR}}^c)$,

$$M^\nu = \frac{1}{2}\begin{pmatrix} 0 & \frac{v}{\sqrt{2}}\mathbf{h}^\nu \\ \frac{v}{\sqrt{2}}\mathbf{h}^{\nu T} & \mathbf{m} \end{pmatrix}, \quad (46)$$

where $\mathbf{h}^\nu = (h^\nu_{ij})$. We analogously denote by boldface the various Yukawa matrices appearing in Eq. (5), such as $\mathbf{h}^E$, $\mathbf{h}^U$, $\mathbf{h}^d$, $\mathbf{f}^u$, $\mathbf{f}^D$. The mass matrix (46) has the usual seesaw texture which generates masses at the sub-eV scale for the left-handed neutrinos assuming, for example, h^ν_{ij} of order one and $\mathbf{m} \sim 10^{14}$ GeV.

As we have already pointed out in Sec. II A, due to the residual $U(1)_\mathcal{G}$ symmetry, the Yukawa Lagrangian in Eq. (5) is not sufficient for giving mass to all charged fermion fields. In order to overcome this problem, we consider the following dimension-5 effective operators,[5] which can emerge from a simple ultraviolet completion shown in Sec. V,

$$-\mathcal{L} \supset \frac{y^e_{is}}{\Lambda}\overline{\Psi_{iL}}\chi^*\rho^* e'_{sR} + \frac{y^u_{am}}{\Lambda}\overline{Q_{aL}}\chi\rho u'_{mR} + \frac{y^d_n}{\Lambda}\overline{Q_{3L}}\chi^*\rho^* d'_{nR} + \text{H.c.} \quad (47)$$

The large mass scale is $\Lambda \gg w$ and the matrices of coefficients y^e_{is}, y^u_{am} and y^d_n have sizes 3×6, 2×5 and 1×4 respectively; they are denoted by $\mathbf{y}^e$, $\mathbf{y}^u$ and $\mathbf{y}^d$ when

[5]Effective operators like these have also been considered recently in a similar model in Ref. [17].

the indices are suppressed. A contraction of the SU(3) antisymmetric tensor ε_{ijk} with the triplet fields in Eq. (47) is implicit. These effective operators break explicitly the $\mathrm{U}(1)_{\mathcal{G}}$ symmetry allowing mass generation for the remaining charged fermion fields that are left massless when only Eq. (5) is considered.

The 6×6 charged lepton mass matrix, in the basis $(e_i, E_i)_{L,R}$, has the form

$$\mathcal{M}^l = \frac{1}{\sqrt{2}}\begin{pmatrix} \epsilon \mathbf{y}^e \\ w\mathbf{h}^E \end{pmatrix} = \begin{pmatrix} M_e & M_{eE} \\ 0_{3\times 3} & M_E \end{pmatrix}, \tag{48}$$

where M_e, M_{eE}, M_E are all 3×3 matrices and the first two matrices are hierarchically suppressed by $\epsilon = vw/\sqrt{2}\Lambda \ll v$. We can choose the lower left block to vanish without loss of generality—and without disrupting the natural hierarchy between the first three rows ($\sim\epsilon$) and the last three ($\sim w$)—by rotating appropriately in e'_{sR} whose rotation matrix is unphysical as all fields are singlets of the gauge group. In this form, M_e already approximately represents the mass matrix for the charged leptons l_α, $\alpha = e, \mu, \tau$, of the SM and M_E represents the mass matrix for the heavy charged leptons $\mathcal{E}_i$, $i = 1, 2, 3$. By further exploring the freedom to rotate ψ_{iL} we could make either M_e or M_E diagonal. There is a mixing among the $l_{\alpha L}$ and $\mathcal{E}_{iL}$ controlled by the entry m_{eE} and has magnitude suppressed by $M_{eE}/M_E \sim \epsilon/w$. In the limit $\epsilon \to 0$ ($\Lambda \to \infty$) the eigenstates l_i become massless as a result of the $\mathrm{U}(1)_{\mathcal{G}}$ symmetry restoration. Therefore, it is natural that the leptons l_α are lighter than $\mathcal{E}_i$.

The above scenario associates the mass of the known charged leptons to the energy scale ϵ which is derived from the electroweak scale v times a suppression factor w/Λ. For example, for $w = 10$ TeV and $\epsilon \sim m_\tau \sim 1$ GeV, we have $\Lambda \sim 10^3$ TeV. Although there is still an unexplained fine tuning of order 10^{-3} for the electron mass relative to the scale ϵ, this situation contrasts with the Standard Model where a tuning of order 10^{-5} relative to the electroweak scale is required.

For the up-type quark mass matrix, in the basis $(u_1, u_2, u_3, U_1, U_2)^T_{L,R}$, we obtain similarly

$$\mathcal{M}^u = \frac{1}{\sqrt{2}}\begin{pmatrix} -\epsilon \mathbf{y}^u \\ v\mathbf{f}^u \\ w\mathbf{h}^U \end{pmatrix} = \begin{pmatrix} M_u & M_{uU} \\ 0_{2\times 3} & M_U \end{pmatrix}, \tag{49}$$

after an appropriate redefinition of u'_{mR}, $m = 1, \ldots, 5$; M_u, M_{uU}, M_U are matrices of sizes 3×3, 3×2, and 2×2, respectively. By also rotating Q_{aL} we can choose $M_U = \mathrm{diag}(m_{\mathcal{U}_1}, m_{\mathcal{U}_2})$ as diagonal, whose values of order w correspond to the heavy quarks $\mathcal{U}_1$, $\mathcal{U}_2$ of charge 2/3. Analogously, M_u corresponds to the mass matrix of the up-type quarks of the SM, (u, c, t). The large separation in energy among the sets of rows in (49) naturally suppresses the mixing between states with hierarchically different masses [44]. The mixing between the heavy quarks $\mathcal{U}_{aL}$ and the SM quarks (left-handed) are controlled by the entry M_{uU} and is at most $M_{uU}/M_U \sim v/w \sim 10^{-2}$ for t_L and at most of order $\epsilon/w \sim 10^{-4}$ for (u_L, c_L). The entries of M_u themselves have a natural hierarchy of $\epsilon/v \sim 10^{-2}$ between the first two rows and the third. By conveniently rotating the right-handed components, we can write the mass matrix for the SM up-type quarks,

$$M_u = \frac{1}{\sqrt{2}}\begin{pmatrix} \epsilon y_i^u \\ v f_i^u \end{pmatrix} = \begin{pmatrix} \tilde{M}_u & \tilde{M}_{ut} \\ 0_{1\times 2} & m_t \end{pmatrix}, \tag{50}$$

where m_t is the top mass of order v and $\tilde{M}_u$, $\tilde{M}_{ut}$ are of order ϵ or smaller. The mass matrix $\tilde{M}_u$ is naturally of order $\epsilon \sim 1$ GeV and gives masses for (u, c). We could have chosen $\tilde{M}_u$ to be diagonal instead of M_U from the rotation on Q_{aL}. The mixing between t_L and (u_L, c_L) is naturally suppressed by $\tilde{M}_{ut}/m_t \sim \epsilon/v \sim 10^{-2}$.

For the down-type quarks, in the basis $(d_1, d_2, d_3, D)^T_{L,R}$, we have the mass matrix

$$\mathcal{M}^d = \frac{1}{\sqrt{2}}\begin{pmatrix} v\mathbf{h}^d \\ \epsilon \mathbf{y}^d \\ w\mathbf{f}^D \end{pmatrix} = \begin{pmatrix} M_d & M_{dD} \\ 0_{1\times 3} & M_D \end{pmatrix}, \tag{51}$$

after appropriate rotation in d'_{nR}, $n = 1, \ldots, 4$. The mass M_D of order w corresponds to a new heavy quark $\mathcal{D}$ while SM quarks (d, s, b) have a mass matrix given approximately by M_d. The mixing between $\mathcal{D}_L$ and the (d_L, s_L, b_L) are naturally suppressed by at least $M_{dD}/M_D \sim v/w \sim 10^{-2}$. In fact, for the down-type quarks, we do not obtain a natural hierarchy between the first two families and the third family. We obtain a natural hierarchy if h^d_{an} is not of order one but suppressed by

$$v h^d_{an} = \epsilon' \bar{h}^d_{an} \sim 6\times 10^{-4} v \bar{h}^d_{an}, \tag{52}$$

with $\bar{h}^d_{an}$ of order one and $\epsilon' \sim m_s \sim 0.1$ GeV. This suppression further decreases the mixing between $\mathcal{D}_L$ and b_L to $\epsilon/w \sim 10^{-4}$ and one order of magnitude smaller for the mixing with (d_L, s_L). We show in the following a possible mechanism responsible for this further suppression. Assuming this hierarchy for the moment, we obtain the mass matrix for (d, s, b):

$$M_d = \frac{1}{\sqrt{2}}\begin{pmatrix} \epsilon' \bar{h}^d_{ai} \\ \epsilon y^d_i \end{pmatrix} = \begin{pmatrix} \tilde{M}_d & \tilde{M}_{db} \\ 0_{2\times 1} & m_b \end{pmatrix}, \tag{53}$$

where we have used appropriate rotations on d'_{iR}, $i = 1, 2, 3$. We can see that m_b is naturally of order $\epsilon \sim 1$ GeV and $\tilde{M}_d$—which yields the masses for (d, s)—is naturally of order $\epsilon' \sim m_s$. The mixing between b_L and (d_L, s_L) is naturally suppressed by $\tilde{M}_{db}/m_b \sim m_s/m_b \sim 0.02$.

The necessary suppression in $\mathbf{h}^d$ could arise if the operator $h^d_{an}\overline{Q_{aL}}\rho^* d'_{nR}$ in Eq. (5) is in fact absent at tree-level but results from an effective higher order operator involving a new singlet scalar φ at a very high energy as

$$\bar{h}^d_{an}\frac{\varphi}{\Lambda'}\overline{Q_{aL}}\rho^* d'_{nR} + \text{H.c.} \to \bar{h}^d_{an}\frac{\langle\varphi\rangle}{\Lambda'}\overline{Q_{aL}}\rho^* d'_{nR} + \text{H.c.} \tag{54}$$

Thus, we would have effectively that $h^d_{an} = \bar{h}^d_{an}\langle\varphi\rangle/\Lambda'$, where $\langle\varphi\rangle/\Lambda' \sim \epsilon'/v \sim 6\times 10^{-4} \ll 1$. The absence of the tree-level term $\overline{Q_{aL}}\rho^* d'_{nR}$ can be arranged by introducing a $\mathbb{Z}_2$ symmetry under which only φ, Q_{aL}, u_{aR}, U_{aR} are odd. One of the up-type quarks, u_{3R}, is kept even so that the top mass is still generated correctly by Eq. (5). The direct interaction terms involving the bilinear forms $\overline{Q_{aL}}u_{3R}$, $\overline{Q_{3L}}u_{aR}$ and $\overline{Q_{3L}}U_{aR}$ will be absent but effectively induced by the replacements $\overline{Q_{aL}} \to \varphi/\Lambda'\overline{Q_{aL}}$, $u_{aR} \to u_{aR}\varphi/\Lambda'$, and $U_{aR} \to U_{aR}\varphi/\Lambda'$, so that the mixing between t_L and the heavy $\mathcal{U}_{aL}$ or the lighter (u_L, c_L) will be further suppressed by $\epsilon'/v \sim 10^{-3}$ compared to the estimates discussed above. This property renders the top quark essentially unmixed with the rest.

Therefore, our minimal mechanism of breaking the 3-3-1 symmetry by the use of just two triplets, together with the $\mathbb{Z}_2$ symmetry above, correctly displays the qualitative hierarchy between the masses for the third family quarks and those of the first two families. Comparing the mass matrices in Eqs. (50) and (53), it is clear that we have enough freedom to obtain the correct masses for the SM quarks and the correct Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix in a quantitative way. We can write

$$\begin{aligned} M_d &= (V^d_L)^\dagger \text{diag}(m_d, m_s, m_b), \\ M_u &= (V^u_L)^\dagger \text{diag}(m_u, m_c, m_t), \end{aligned} \tag{55}$$

already discarding the unobservable rotation matrices for the right-handed quarks. The CKM matrix is fixed by

$$V_{\text{CKM}} = (V^u_L)^\dagger V^d_L, \tag{56}$$

so that V^d_L can be considered free, while V^u_L is fixed by V_{CKM} and V^d_L. From the discussion above, V^u_L is essentially block-diagonal with the third family practically decoupled from the rest. Another possibility to naturally suppress the coupling h^d_{an} would be to implement the Froggatt-Nielsen mechanism [45] (see also [46] for a proposal along this line in a 3-3-1 model). We leave this question for future investigations.

We briefly comment on the model with $\beta = -1/\sqrt{3}$ where the heavy quark content is inverted—we would have two heavy quarks D_a of charge $-1/3$ and just one quark U of charge $2/3$—and the heavy charged leptons E_i are replaced by neutral leptons N_i that could participate in the mass generation mechanism for light neutrinos. The natural hierarchy in the up-type and down-type quark sectors would be very similar and the implementation of the $\mathbb{Z}_2$ symmetry is analogous.

IV. SUPPRESSED FLAVOR CHANGING CURRENTS

We consider first the FCNC interactions mediated by scalars. To that end it is instructive to consider first the breaking of $SU(3)_L \otimes U(1)_X$ to the SM gauge group by $\langle\chi\rangle$ and rewrite

$$\chi = \begin{pmatrix} -\phi_1 \\ \chi^0_3 \end{pmatrix}, \qquad \rho = \begin{pmatrix} \tilde{\phi}_2 \\ \rho^-_3 \end{pmatrix}, \tag{57}$$

where ϕ_1, ϕ_2 are $SU(2)_L$ doublets of $Y = 1/2$. At this stage, only $\sqrt{2}\text{Re}\chi^0_3$ acquires a vev $w \sim 10$ TeV and $\sqrt{2}\text{Re}(\chi^0_3)$ and ρ^-_3 will be heavy SM singlet scalars of $Y = 0$ and $Y = -1$, respectively. The scalar $\sqrt{2}\text{Im}\chi^0_3$ and the doublet ϕ_1 will be absent in the unitary gauge because they will be absorbed by the gauge bosons Z'_μ and $(V^+_\mu, V^0_\mu)^T$.[6] See Sec. II B for their composition after EWSB. When taking into account the effective operators in (47), it is convenient to write

$$\frac{1}{\Lambda}\chi^* \times \rho^* = \begin{pmatrix} \phi_3 \\ \frac{\phi_1^\dagger\phi_2}{\Lambda} \end{pmatrix}, \tag{58}$$

where

$$\phi_3 = \frac{1}{\Lambda}(\chi^{0*}_3\phi_2 - \rho^+_3\tilde{\phi}_1) \sim \frac{\text{Re}\chi^0_3}{\Lambda}\phi_2, \tag{59}$$

is a SM effective Higgs doublet. The dominant contribution coming from (47) at this stage will be

$$\phi_3 \to \frac{\langle\chi^0_3\rangle}{\Lambda}\phi_2 = \frac{\epsilon}{v}\phi_2 \sim 10^{-2}\phi_2. \tag{60}$$

We also separate the quark triplets in Eq. (3) and lepton triplets in Eq. (2) into SM doublets and singlets as

$$Q_{aL} = \begin{pmatrix} i\sigma_2 q_{aL} \\ U_{aL} \end{pmatrix}, \qquad Q_{3L} = \begin{pmatrix} q_{3L} \\ D_L \end{pmatrix}, \qquad \psi_{iL} = \begin{pmatrix} l_{iL} \\ E_{iL} \end{pmatrix}, \tag{61}$$

[6] The charged component of ϕ_1 will have a small admixture with ρ^+_3 after Electroweak symmetry breaking (EWSB); see Eq. (22).

where q_{aL} is the usual quark doublet of family $a = 1, 2$, σ_2 is one of the Pauli matrices, and l_{iL} are the usual lepton doublets.

At this stage of breaking, the model is equivalent to the SM with additional heavy singlet vectorlike quarks (VLQ) D, $U_{1,2}$, singlet vectorlike leptons E_i and heavy singlet scalars $\sqrt{2}\mathrm{Re}(\chi_3^0)$, $\rho_3^{\pm}$, with the additional heavy gauge bosons. The Yukawa interactions in (5) and (47) that only involve the light doublet ϕ_2 are

$$\begin{aligned}-\mathcal{L}_0 = {}& \bar{q}_{aL}\phi_2\frac{\sqrt{2}}{\epsilon'}[(M_d)_{ai}d_{iR} + (M_{dD})_a D_R] \\ &+ \bar{q}_{3L}\phi_3\frac{\sqrt{2}}{\epsilon}[m_b d_{3R} + (M_{dD})_3 D_R] \\ &+ \bar{q}_{aL}\tilde{\phi}_3\frac{\sqrt{2}}{\epsilon}[(M_u)_{ai}u_{iR} + (M_{uU})_{ab}U_{bR})] \\ &+ \bar{q}_{3L}\tilde{\phi}_2\frac{\sqrt{2}}{v}[m_t u_{3R} + (M_{uU})_{3b}U_{bR}] + \mathrm{H.c.} \end{aligned} \tag{62}$$

See the Appendix for other interactions involving heavy particles. Note that $d_{3R} \approx b_R$, $u_{3R} \approx t_R$ and $q_{3L} \approx (t_L, b_L)^{\mathsf{T}}$ are almost the mass eigenstates except for suppressed mixing with the lighter quarks or the heavier $\mathcal{D}$, $\mathcal{U}_a$. With only one effective Higgs doublet, there is natural flavor conservation and no flavor changing neutral interactions mediated by scalars [47], except for the ones induced by the small mixing between SM and heavy quarks [44]. This contrasts with the usual mechanism for breaking the 3-3-1 symmetry involving three $\mathrm{SU}(3)_L$ triplets: usually there are two light Higgs doublets which induce suppressed but nonvanishing neutral flavor changing interactions at tree level [48].

We can consider now the currents coupling with the gauge bosons. There are two types of FCNC interactions for SM fermions: (i) the ones coming from the mixing between the third family of quarks and the first two families and (ii) the ones coming from the mixing between heavy fermions and the SM fermions. The first type (i) inevitably appears in all 3-3-1 models because one of the quark families is treated differently by the $\mathrm{SU}(3)_L$. Treating the third family differently is the only option if we want to avoid unrealistic flavor changing contributions in the mass differences of the K^0, D^0 and B^0 systems for a Z_2^μ of mass no larger than a few TeV [48]. Sufficiently suppressed mixing of the third family with the lighter ones thus leads to experimentally allowed flavor changing contributions. In our model, however, because of the different mass generation mechanism, such a mixing is already *naturally* suppressed at least at the level of $\epsilon/v \sim 10^{-2}$ for the up-type quarks (further suppressed with the $\mathbb{Z}_2$ implementation) and at least $m_s/m_b \sim 10^{-2}$ for the down-type quarks. The flavor changing interactions of type (ii) are well known to be naturally suppressed by the hierarchy of SM fermions and heavy fermions [44]. We discuss these suppressed interactions in more detail in the following.

The fermionic currents couple with a vector boson V^μ as

$$\mathcal{L} \supset -V_\mu J_V^\mu. \tag{63}$$

For the neutral gauge boson $V = Z'$, we have

$$J_{Z'}^\mu = g_{Z'}\bar{\Psi}\gamma^\mu[s_W^2 Y - \sqrt{3}c_W^2 T_8]\Psi, \tag{64}$$

where $g_{Z'} \equiv \frac{g}{\sqrt{3}c_W\sqrt{1-4s_W^2/3}}$ and Ψ is the collection of all the fermions in the symmetry basis. The flavor changing piece can be extracted as

$$\begin{aligned} g_{Z'}^{-1}J_{Z'}^\mu|_{\mathrm{FCNC}} = {}& (-c_W^2)\bar{u}_{3L}\gamma^\mu u_{3L} + (-c_W^2)\bar{d}_{3L}\gamma^\mu d_{3L} \\ &+ \left(-\frac{3}{2} + 2s_W^2\right)\bar{U}_{aL}\gamma^\mu U_{aL} \\ &+ \left(\frac{1}{2} - s_W^2\right)\bar{D}_L\gamma^\mu D_L \\ &+ \left(c_{2W} + \frac{1}{2}\right)\bar{E}_{iL}\gamma^\mu E_{iL}, \end{aligned} \tag{65}$$

where $c_{2W} \equiv \cos 2\theta_W$, and we have subtracted the flavor universal part common to the first two families of quarks and three family of leptons. Small mixing terms appear when we go to the basis of physical fields of definite masses. We can clearly see in (65) the two types of FCNC discussed above: the first line induces interactions of type (i) while the rest induces the type (ii) currents. For type (i), only the mixing between the third family and the first two families will be observable. This mixing is naturally suppressed by $\epsilon/v \sim m_s/m_b \sim 10^{-2}$ in our model which is typically below the current limits coming from meson mixing [49].

As the interaction with the Z boson is identical to the ones in the SM for the usual quarks and leptons, FCNCs are only of type (ii):

$$g_Z^{-1}J_Z^\mu|_{\mathrm{FCNC}} = -\frac{1}{2}\bar{U}_{aL}\gamma^\mu U_{aL} + \frac{1}{2}\bar{D}_L\gamma^\mu D_L + \frac{1}{2}\bar{E}_{iL}\gamma^\mu E_{iL}, \tag{66}$$

where $g_Z \equiv g/c_W$. We have again subtracted the family universal contributions from the first families of quarks and leptons. The coupling with W is the same as in the SM in the symmetry basis and small nonunitary effects appear through mixing between heavy and SM fermions. Flavor changing interactions are constrained by the search for singlet VLQs at the LHC [50] and by indirect constraints coming from precision electroweak observables and Large Electron-Positron Collider (LEP) [51]. The former constrains the masses to be above around 1 TeV and the latter constrains the mixing angle between the heavy quarks and the third family to be less than 0.04 for the down-type quarks and 0.14 for the up-type quarks for heavy quarks of 1 TeV. So our heavy quarks of masses at the scale $w \sim 10$ TeV and mixing angle of less than 10^{-2} are not currently observable. If the heavy quark masses are lowered to few TeV, and the $\mathbb{Z}_2$ that decouples the top is present, the dominant channel for $\mathcal{U}_a$ will be $\mathcal{U}_a \to Wb + X$, as $\mathcal{U}_a \to ht + X$ and $\mathcal{U}_a \to Zt + X$ will be negligible. For $\mathcal{D}$, the channels $\mathcal{D} \to hb + X$ and $\mathcal{D} \to Wt + X$ are similarly important. The constraints for singlet vectorlike leptons are much more relaxed.

For completeness, we also collect the interactions with the heavy gauge bosons V^0, V^+:

$$-\mathcal{L}_V = \frac{g}{\sqrt{2}}\left[-\bar{q}_{aL}\gamma^\mu \begin{pmatrix} V_\mu^0 \\ -V_\mu^- \end{pmatrix} U_{aL} + \bar{q}_{3L}\gamma^\mu \begin{pmatrix} V_\mu^+ \\ V_\mu^0 \end{pmatrix} D_L + \bar{l}_{iL}\gamma^\mu \begin{pmatrix} V_\mu^+ \\ V_\mu^0 \end{pmatrix} E_{iL}\right] + \text{H.c.}, \tag{67}$$

where the heavy gauge bosons $(V_\mu^+, V_\mu^0)^\mathsf{T}$ have the same gauge quantum numbers as the SM Higgs doublet. These gauge bosons lie at the scale w and interactions with two SM fermions are suppressed by the heavy-light mixing.

V. AN ULTRAVIOLET COMPLETION OF THE MODEL

We show in this section a simple ultraviolet completion of the model allowing for the generation of the effective operators in Eq. (47). In order to achieve that, we add a scalar field $\eta \sim (\mathbf{1}, \mathbf{3}, 1/3)$ which transforms in the same way as χ under $\mathrm{SU}(3)_C \otimes \mathrm{SU}(3)_L \otimes \mathrm{U}(1)_X$. We assume that η has a mass $M \gg w, v$ much larger than the rest of the other fields in the model. Thus, as we describe below, η can be integrated out so that the remaining effective theory is exactly the model studied above with the dimension-5 effective operators given in (47).

With the introduction of η, the total scalar potential is

$$V_T(\eta, \chi, \rho) = V(\chi, \rho) + V_\eta, \tag{68}$$

where $V(\chi, \rho)$ is given in Eq. (6) and

$$\begin{aligned} V_\eta = {}& M^2\eta^\dagger\eta + (\lambda_5(\eta^\dagger\eta) + \lambda_6(\rho^\dagger\rho) + \lambda_7(\chi^\dagger\chi))(\eta^\dagger\eta) \\ &+ \lambda_8(\eta^\dagger\rho)(\rho^\dagger\eta) + \lambda_9(\eta^\dagger\chi)(\chi^\dagger\eta) \\ &- [f\eta\rho\chi - \lambda_{10}(\eta^\dagger\rho)(\rho^\dagger\chi) - (\lambda_{11}\eta^\dagger\chi + \lambda_{12}\eta^\dagger\eta \\ &+ \lambda_{13}\chi^\dagger\chi + \lambda_{14}\rho^\dagger\rho)(\eta^\dagger\chi) + \text{H.c.}], \end{aligned} \tag{69}$$

where $M^2 > 0$ is the quadratic mass for η; with the coupling constant f, which we take as being real, having dimension of mass; and the λ's ($< 4\pi$) are the usual scalar field perturbative self-interaction coupling constants. Also, we consider a basis (χ, η) in which the bilinear terms in these fields are diagonal. In fact, bilinear terms, such as $\mu_3^2\eta^\dagger\chi$, can be eliminated through a rotation to the diagonal basis, implying effectively a change on the original quadratic mass parameters along with a redefinition of the quartic coupling constants.

We can see that the approximate conservation of the global charge $\mathcal{G}$ in this UV completion is guaranteed by the fact that the breaking is induced solely by a soft breaking of the PQ symmetry through the f term in (69). Therefore the breaking effects are all proportional to the breaking parameter f even if we consider radiative corrections, and this fact justifies the approximate conservation of $\mathcal{G}$ at low energies.

The Yukawa interactions involving η are similar to those for χ in Eq. (5), adding to such an equation the terms

$$\mathcal{L}_Y \supset y_{is}^e \overline{\Psi_{iL}}\eta e'_{sR} + y_{am}^u \overline{Q_{aL}}\eta^* u'_{mR} + y_n^d \overline{Q_{3L}}\eta d'_{nR} + \text{H.c.} \tag{70}$$

Assuming $M \gg |f| \gtrsim w$, at low energies η is effectively given by

$$\eta \approx \frac{f}{M^2}\rho^*\chi^* + \cdots, \tag{71}$$

where the ellipsis stands for operators which are even more suppressed by M.[7] Replacing this last expression for η in Eq. (70), we get the effective operators in Eq. (47) with the identification $\Lambda = M^2/f$. Moreover, there exist corrections to tree level parameters that shift the couplings $\lambda_3 \to \lambda_3 - |f|^2/M^2$ and $\lambda_4 \to \lambda_4 + |f|^2/M^2$. As an example, the value $\Lambda = 983$ TeV can be achieved with $M \approx 10^5$ GeV and $f \approx 10^4$ GeV. We see that for $|\lambda_{3,4}|$ of order unity, the correction $|f|^2/M^2 \sim 10^{-2}$ does not have a significant impact on those couplings.

VI. PHENOMENOLOGY OF THE Z' BOSON

In this section, we present some results involving the new neutral vector boson within the context of the LHC at the 14 TeV energy regime. By the reason that the mixing angle θ in Eq. (44) between Z and Z' is small for $w \gg v$, we have $Z_2 \approx Z'$. Thus, we consider in the following analysis the new vector boson as being Z' and its couplings with fermions. Constraints on the Z' mass coming from FCNCs can be strongly dependent on the specific model of choice. In particular, for 3-3-1 models with $\beta = \pm 1/\sqrt{3}$, it has been shown that by choosing either the first or the third quark family to transform differently from the others leads to different constraints on the w scale and, consequently, on

[7]There are other terms of the same order in $1/M^2$ that correct the contribution of χ but only (71) leads to a vev in a direction orthogonal to $\langle\chi\rangle$ and $\langle\rho\rangle$.

the new gauge boson masses [17]. For other versions, a lower-bound of 3 TeV for the Z' mass has been found [52].

In addition, previous studies concerning the Z' branching ratios for the $\beta = \pm\sqrt{3}$ versions have identified a leptophobic character of such a neutral gauge boson [53]. In our case, within a scenario where the exotic masses are just above 1 TeV, the Z' branching ratios are divided into $\mathrm{Br}[Z' \to \nu\bar{\nu}] \simeq 45\%$, $\mathrm{Br}[Z' \to \ell\bar{\ell}] \simeq 13\%$ and $\mathrm{Br}[Z' \to q\bar{q}] \simeq 42\%$. Thus, when we compare the leptonic Z' branching ratio with the SM $\mathrm{Br}[Z \to \ell\bar{\ell}] \simeq 3\%$, we can conclude that the search for the new gauge boson can be accessible via a clean dilepton signal at the LHC. Moreover, from the relation among the gauge boson masses, it is clear that our Z' can only decay into fermions and scalars, in contrast with the leptophobic versions where the channels involving the new $\mathrm{SU}(3)_L$ gauge bosons are present. Finally, by calculating the Z' width, we find that $\Gamma_{Z'}$ is around $5\% M_{Z'}$.

Then, by considering the possibility of the $\mathrm{SU}(3)_L$ breaking scale being at the $\mathcal{O}(\mathrm{TeV})$, leading to a mass scale for the new gauge bosons of a few TeVs, we explore the production of a muon pair through the decay of the heavy Z'. It is clear that, from the experimental point of view, the new neutral gauge boson can be observed in the invariant mass formed by the dilepton mass spectrum. The peak observed in the invariant mass distribution for the final particles, over a smooth SM background, represents the evidence for new physics. Thus, in general, the experimental analysis searches for narrow resonances where the experimental resolution is the dominant contribution to the observable width of a peak structure appearing over a SM background. In this approach, theoretical cross section predictions for specific models are usually calculated in the narrow width approximation. Obviously, when the width is wide, the resonance appears as a broad shape and can be almost flat around the Z' pole.

Thus, within this narrow width approximation, we show below the invariant mass and transverse momentum p_T distributions of the emerging leptons in the processes $p + p \to \mu^+ + \mu^- + X$ at 14 TeV, involving the Z' of this new 3-3-1 version. We leave for a future work the study of the effects of a Z' with wide width, like the one predicted in the so-called minimal version.

To carry out our analysis, we consider the general Lagrangian for the neutral currents involving Z and Z' contributions,

$$\mathcal{L}^{NC} = -\frac{g}{2\cos\theta_W}\sum_f [\bar{f}\gamma^\mu(g_V + g_A\gamma^5)fZ_\mu + \bar{f}\gamma^\mu(g'_V + g'_A\gamma^5)fZ'_\mu], \tag{72}$$

where f stands for leptons and quarks, g is the weak coupling constant, and g_V, g_A, g'_V and g'_A, are the SM and 3-3-1 couplings which are presented in the Table II, where we take the approximation $v/w \ll 1$ and assume no flavor mixing. Below the electroweak scale, the phenomenology predicted by the new model involving γ and Z coincides with the SM one.

TABLE II. The vector and axial couplings of Z and Z' to leptons (e, μ and τ) and quarks (u and d) in the 3-3-1 models. θ_W is the Weinberg angle.

	g_V	g_A	g'_V	g'_A
$Z\bar{l}l/Z'\bar{l}l$	$-\frac{1}{2} + 2\sin^2\theta_W$	$-\frac{1}{2}$	$-\frac{1+2\sin^2\theta_W}{2\sqrt{3-4\sin^2\theta_W}}$	$-\frac{1-2\sin^2\theta_W}{2\sqrt{3-4\sin^2\theta_W}}$
$Z\bar{u}u/Z'\bar{u}u$	$\frac{1}{2} - \frac{4\sin^2\theta_W}{3}$	$\frac{1}{2}$	$\frac{3+2\sin^2\theta_W}{6\sqrt{3-4\sin^2\theta_W}}$	$\frac{1-2\sin^2\theta_W}{2\sqrt{3-4\sin^2\theta_W}}$
$Z\bar{d}d/Z'\bar{d}d$	$-\frac{1}{2} + \frac{2\sin^2\theta_W}{3}$	$-\frac{1}{2}$	$\frac{\sqrt{3-4\sin^2\theta_W}}{6}$	$\frac{1}{6\sqrt{3-4\sin^2\theta_W}}$

By following previous studies on Z' concerning strategies for the identification of this particle on the muon channel [54–56], as well as the last ATLAS report [57], we have applied some cuts in order to obtain clear distributions for the invariant masses and transverse momentum of the final muons. In agreement with the ATLAS detector performance, the cuts adopted for the pseudorapidity and the transverse momentum of the muons are: $|\eta| < 2.5$ and $p_T > 30$ GeV. For the invariant mass of the muon pair, we have used a strong cut ($M_{\mu\mu} > 1000$ GeV) in order to suppress the SM background.

In our simulations, we have made use of the CompHep [58] and the MadAnalysis [59] packages and adopted the CTEQ6L [60] parton distribution functions set, evaluated at the $\sqrt{\hat{s}}$ factorization/renormalization scale, i.e., the center-of-mass energy at the parton level.

Upon assuming 3, 4 and 5 TeV for the Z' mass, we observe the resonance peaks around the respective masses in the invariant mass distribution. If we consider two values for the projected LHC integrated luminosity ($\mathcal{L} = 100\ \mathrm{fb}^{-1}$ and $\mathcal{L} = 300\ \mathrm{fb}^{-1}$) at $\sqrt{s} = 14$ TeV, we obtain the number of events as shown in Fig. 1. As the width of the heavy boson satisfies the relation $\Gamma_{Z'} \sim 5\% M_{Z'}$, our results are in accordance with Ref. [56].

In order to identify the new gauge boson, we also consider the muon p_T distribution, where two peaks are expected to appear, corresponding to one half of the resonance masses (M_Z, $M_{Z'}$), but, in this case, due to the invariant mass cut adopted, the first peak moves justly to one half of this cut (~500 GeV). By applying an additional cut in the muon transverse momentum, $p_T > 500$ GeV, we obtain the distributions shown in the Fig. 2. We can identify the peak around the $M_{Z'}/2$ values, with $M_{Z'} = 3$, 4 and 5 TeV, respectively. It is clear that, for $M_{Z'} = 3$ and 4 TeV, the peak is well defined, and for greater masses, the peak is smooth like the SM background. Then, with a stronger cut on the invariant mass or on the muon p_T, we might obtain a clearer peak for higher masses in the p_T distribution.

Therefore, as claimed in [54,56], we can use the transverse momentum distribution of the final muons as an

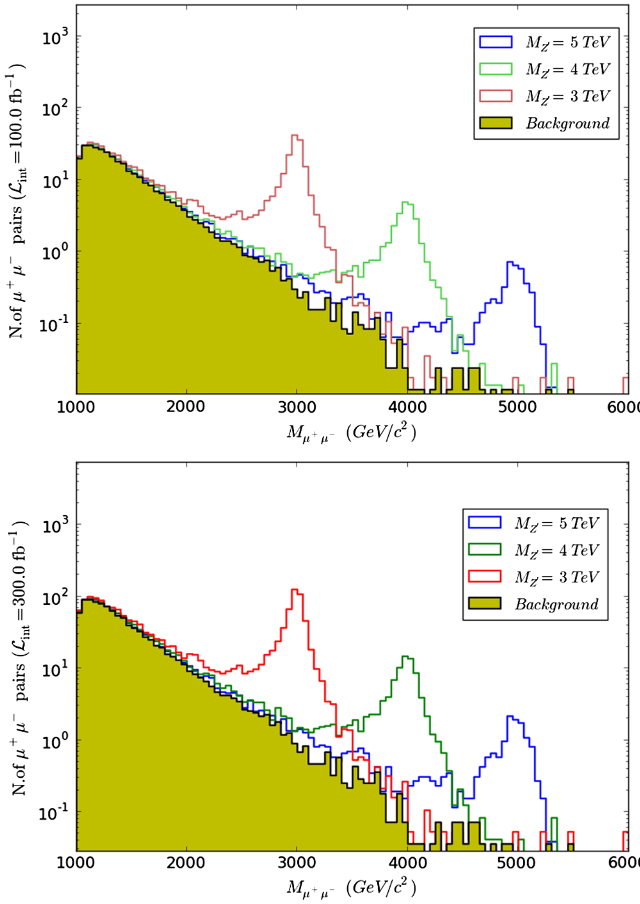

FIG. 1. Number of $\mu^+\mu^-$ pairs as a function of three representative Z' masses at the 14 TeV LHC with $\mathcal{L} = 100$ fb^{-1} (upper panel) and $\mathcal{L} = 300$ fb^{-1} (lower panel).

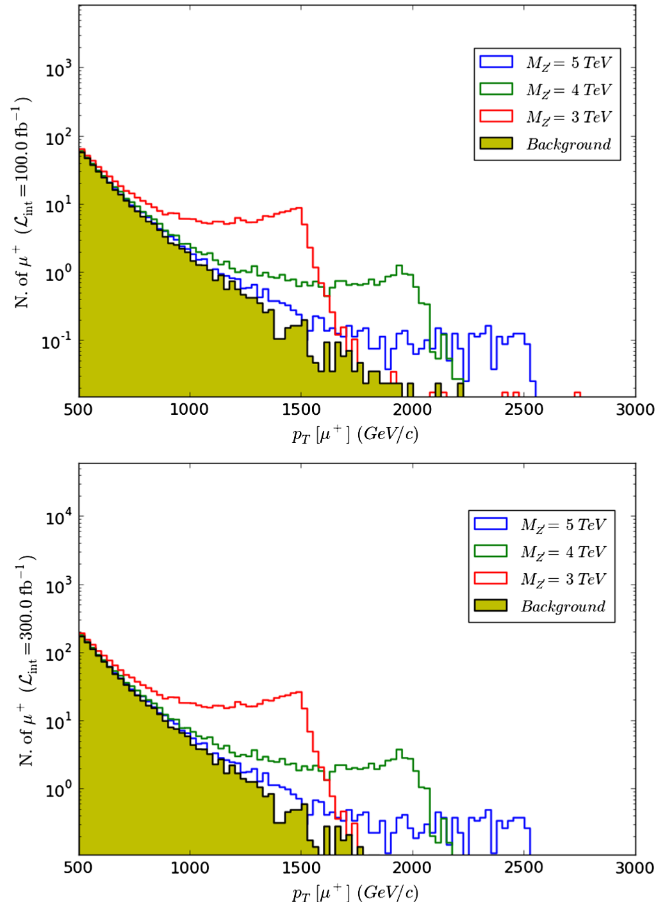

FIG. 2. Number of μ^+ as a function of the p_T of the emerging μ at the 14 TeV LHC with $\mathcal{L} = 100$ fb^{-1} (upper panel) and $\mathcal{L} = 300$ fb^{-1} (lower panel).

additional tool to distinguish the signal coming from the Z' from the SM background. Thus, by a simple analysis of the invariant masses and transverse momentum of the final μ, we get clear signals for the Z'. Moreover, by considering the projected luminosities for the LHC run-II, we obtain a considerable number of events revealing the existence of the new gauge boson. In a recent work [61], a phenomenological analysis of the Z' has been performed for the version considered in this paper. The authors have suggested that the possibility of detecting the new particle can be achieved by considering just the production of a pair of leptons. From their analysis using, for example, the Forward-Backward asymmetry, they have concluded that at the LHC at 14 TeV, it is possible to identify the Z' boson. Our strategy is a little different from the one used by those authors. We use the invariant mass distributions and the transverse momentum of the final leptons, in order to distinguish the signal from the SM background. In any case, our conclusions are similar regarding the possibility of discovering the Z' in the run II of the LHC.

It is beyond the scope of our work to make a detailed analysis of the final states, including $Z - Z'$ interferences, detector efficiencies, hadronization, etc. However, based on our results, it is not hard to establish the existence or to exclude the Z' predicted by this model. Finally, a complete analysis involving the Z' predicted by different versions of the 3-3-1 model within the next stage of the LHC energy is mandatory, but we postpone this study to a future work.

VII. CONCLUSIONS

In this work we have presented a version of the 3-3-1 model, defined by $\beta = 1/\sqrt{3}$ in the electric charge operator in Eq. (1), which at low energies contains only two scalar triplets in order to achieve the correct breakdown of gauge symmetries. Eight out of the twelve degrees of freedom contained in the scalar triplets are absorbed in the longitudinal components of the vector bosons Z, $W^{\pm}$, Z', $V^{\pm}$, V^0, $V^{0\dagger}$. This leaves three spin-0 bosons in the particle spectrum, two neutral CP even scalars h, H, and a charged scalar $\varphi^{\pm}$. The neutral scalar, h, is identified with the discovered Higgs boson, and gets its mass at the scale v, related to the $SU(2)_L \otimes U(1)_Y \to U(1)_Q$ symmetry breakdown. Both the neutral, H, and the charged, $\varphi^{\pm}$, scalars are supposedly heavier, since they get their masses at the scale w, associated with the $SU(3)_L \otimes U(1)_X \to SU(2)_L \otimes U(1)_Y$ symmetry breakdown. In comparison to other Standard Model extensions, such as the Two Higgs Doublet Model for example, our construction has a smaller number of extra scalars at the TeV scale because no CP odd neutral state is part of the spectrum. For the same reason, FCNC mediated by scalars are very much suppressed and absent in the limit of no mixing between SM fermions and heavy fermions.

The model has three extra charged leptons, $\mathcal{E}_i$, two up-type quarks $\mathcal{U}_a$, and one down-type quark $\mathcal{D}$, beyond the Standard Model fermion content. Although the model with just two scalar triplets has a consistent pattern of gauge symmetry breakdown to the electromagnetic factor $U(1)_Q$, some of the standard fermionic fields remain massless due to a residual global $U(1)_\mathcal{G}$ symmetry which, as we observed in Eq. (12), involves diagonal generators of spontaneously broken gauge symmetries plus a sort of Peccei-Quinn symmetry. This $U(1)_\mathcal{G}$ symmetry seems to be a common feature of 3-3-1 models with just two scalar triplets and has also been identified in another version of the model [42]. To overcome this problem, we have introduced a heavy scalar triplet with mass $M \gg w$, which is integrated out from the low energy theory leaving it with effective operators breaking $U(1)_\mathcal{G}$ explicitly, completing the mass generation mechanism for the fermions. As we have shown, the effective operators furnish a less fine-tuned mass generation for leptons and up-type quarks compared to the Standard Model. Such a mechanism, however, does not work as naturally for the standard down-quark mass hierarchy and a solution based on an additional $\mathbb{Z}_2$ symmetry has also been provided. Natural hierarchies between the heavy quarks and the third family quarks, and between the third family and the lighter two families, arise, and this feature naturally suppresses the mixing between them leading to suppressed FCNC interactions.

From the phenomenological standpoint, we have explored the possibility to discover the predicted Z' by considering the leptonic decay channel within the LHC energy regime. By making a simple analysis involving the invariant masses and transverse momentum of the final muons, and by selecting appropriate cuts for the final states, we have concluded that clear signals can reveal the presence of the new neutral gauge boson. If we take the projected integrated luminosities for the next LHC phase, we find a considerable number of events for processes involving Z' and the final muons, which could confirm one of the predictions of the model. Moreover, other potential tests involve the pair and single production of the new leptons and quarks, in addition to the $V^\pm$, V^0, $V^{0\dagger}$ vector bosons. For the minimal scalar sector, containing only an extra Higgs and a charged scalar, the production of H via gluon fusion, i.e. $gg \to H$, and the analysis of the final states $b\bar{b}b\bar{b}$, $b\bar{b}\tau\tau$ and $b\bar{b}\gamma\gamma$ represent an excellent prospect for the discovery or exclusion of the new neutral scalar state. On the other hand, the associated productions of the charged Higgs with a top quark and with a W boson, via the partonic processes $bg \to tH^-$ and $b\bar{b} \to H^- W^+$, can be also tested at the LHC. Thus, considering both the theoretical and the phenomenological aspects presented, this new model is surely worth our attention in further studies.

ACKNOWLEDGMENTS

This research was partially supported by the Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq), by Grants No. 306636/2016-6 (A. G. D.) and No. 308578/2016-3 (C. C. N.). Financial support by Fundação de Amparo à Pesquisa do Estado de São Paulo (FAPESP) is also acknowledged under the Grants No. 2013/22079-8 (A. G. D. and C. C. N.) and No. 2014/19164-6 (C. C. N.). J. L. and R. O. would like to thank CAPES (Coordenação de Aperfeiçoamento de Pessoal de Nível Superior), and W. C. V. thanks UFABC, for the financial support.

APPENDIX: HIGHER ORDER OPERATORS

We collect here the terms that we have omitted in (62):

$$\begin{aligned}-\mathcal{L}_0 \supset & -\bar{q}_{3L}\phi_1 \frac{\sqrt{2}}{w} M_D D_R + \bar{q}_{aL}\tilde{\phi}_1 \frac{\sqrt{2}}{w} M_{U_a} U_{aR} \\ & + \bar{D}_L \chi_3^0 \frac{\sqrt{2}}{w} M_D D_R + \bar{U}_{aL}\chi_3^{0*} \frac{\sqrt{2}}{w} M_{U_a} U_{aR} \\ & - \frac{\phi_2^\dagger \phi_1}{\Lambda} \bar{U}_{aL} \frac{\sqrt{2}}{w} [(M_u)_{ai} u_{iR} + (M_{uU})_{ab} U_{bR})] \\ & + \frac{\phi_1^\dagger \phi_2}{\Lambda} \bar{D}_L \frac{\sqrt{2}}{w} [m_b d_{3R} + (M_{dD})_3 D_R)] + \text{H.c.}\end{aligned} \tag{A1}$$

$$\begin{aligned}-\mathcal{L}_\pm \supset & \bar{D}_L \frac{\sqrt{2}}{v} [m_t u_{3R} + (M_{uU})_{3b} U_{bR}] \rho_3^- \\ & + \bar{U}_{aL} \frac{\sqrt{2}}{\epsilon'} [(M_d)_{ai} d_{iR} + (M_{dD})_a D_R)] \rho_3^+ + \text{H.c.}\end{aligned} \tag{A2}$$

Note that there are also effective Yukawa interactions involving $\rho_3^\pm$ in (62) coming from the terms within ϕ_3.

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