Radiatively Induced Type II seesaw and Vector-like 5/3 Charge Quarks

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Abstract

Understanding small neutrino masses in type II seesaw models with TeV scale SM triplet Higgs bosons requires that its coupling with the standard model Higgs doublet $H$ be “dialed” down to be order eV to KeV, which is a fine-tuning by a factor of $10^{-11} - 10^{-8}$ with respect to the weak scale. We present a SUSY extension of the type II seesaw model where this dimensionful small coupling is radiatively induced, thus making its smallness natural. This model has an exotic vector-like quark doublet which contains a quark $X$ with electric charge $5/3$ and a top partner $t'$. We discuss in details the phenomenology of the model paying special attention to the consequences of the interactions of the the exotic heavy quarks and the scalars of the model. Implications for neutrinoless double beta decay and for the LHC experiments are discussed in detail. Remarkably, in this model both the seesaw triplet and the heavy quarks can manifest at colliders in a host of different signatures, including some that significantly differ from those of the minimal models. Depending on the choice of the hierarchy of couplings, the decay of the heavy quarks and of the seesaw triplet may be subject to bounds that can be tighter or looser than the bounds from standard LHC searches. Furthermore we point out a new short-distance contribution to neutrinoless double beta decay mediated by the simultaneous propagation of the type II triplet and exotic fermions. Remarkably this contribution to the neutrinoless double beta decay is parametrically quite independent from the scale of the generated neutrino mass.
Contents

1 Introduction 2

2 5/3 charge quark and “natural” type II seesaw 4
2.1 Fine tuning problem in type II seesaw 4
2.2 Vector-like quark extension and alleviation of fine-tuning 5
2.3 Quark masses and mixings to $X_{5/3}$ and $t'$ 7

3 Decays of the $\frac{5}{3}$ charge quark and the seesaw triplet 8
3.1 SUSY modes 13

4 New contributions to neutrinoless double beta decay from 5/3 charge quarks 16

5 Other implications of the model 18

6 Conclusions 19

1 Introduction

Seesaw mechanism provides a simple way to understand the smallness of the neutrino mass in gauge theories. It relates the neutrino mass to the inverse of a high mass for new fields added to the standard model e.g. gauge singlet fermion (in type I) [1] or SM triplet Higgs boson in type II seesaw [2] case. Their mass terms being SM singlets are unconstrained by the standard model (SM) gauge symmetry and are allowed to be as high as needed. For order one couplings of leptons to these new fields, this mass scale has to be in the range of $10^{14}$ GeV to get neutrino masses in the eV to sub-eV range. As a result, the presence of new fields and hence these simple seesaw mechanisms becomes hard to test by any low energy or collider experiments except that with additional assumptions like supersymmetry, they may lead to indirect observable signals in processes such as lepton flavor violation, depending on the seesaw scale. With LHC searches for TeV scale new particles in full swing, the possibility that the new fermions and bosons connected with seesaw type I or II may have masses in the TeV range has attracted considerable interest [3]. However, in such a situation, some parameters in the seesaw relation have to be dialed down to rather small values and tend to look less natural. It is therefore important to search for theoretical settings where the input parameters to the seesaw formula arises only as higher order effects so that a TeV scale would look natural and provide a stronger case for searches of the seesaw particles at LHC. In this article we present a model where a key lepton number violating parameter that goes into the type II seesaw formula for neutrino masses is naturally vanishing at the tree level and arises only at the one loop level so that it allows the associated seesaw bosons to have masses in the TeV range in a natural manner [4]. The literature on TeV scale models where neutrino mass arises at the one or two loop levels, outside the seesaw framework, is
enormous \cite{5} and there have also been extensive studies of higher dimensional operators and how they contribute to TeV scale models for eV scale neutrino masses \cite{6}.

Our model has several testable phenomenological consequences. A distinctive feature of the model with respect to ordinary seesaw models is the introduction of a new quark with electric charge $5/3$ (denoted by $X_{5/3}$) as the weak isospin partner of a top partner $t'$. The $(X_{5/3}, t')$ form a vector-like SM doublet, which has a gauge-invariant mass term. It is quite natural in the model for the $X_{5/3}$ particle to have TeV to sub-TeV masses so that it can be searched at the CERN LHC.

The $t'$ has electric charge $2/3$ and after symmetry breaking it mixes with the $u, c, t$ quarks of the SM. As it happens in standard scenarios of heavy quarks this mixing induces a decay of $t'$ into SM weak gauge bosons and SM quarks. Through this same mixing, the $t'$ changes the couplings of the SM quarks to the weak vectors bosons. Both aspects of the $t'$ phenomenology have been extensively studied \cite{7}.

The $X_{5/3}$ particle has been discussed in other contexts \cite{7} such as little Higgs or composite Higgs models. Interestingly, here we are led to the same kind of particles from neutrino mass point of view. While the dominant decay mode of the $X_{5/3}$ particle can still be

\[ X_{5/3} \to tW, \]

as in most models where this particle is introduced, our type II seesaw model endows it with a significant branching ratio to other states of the Higgs sector, such as

\[ X_{5/3} \to \Delta^{++} d_i, \]

which gives rise to multi-leptons final states such as $b\bar{b}\ell^+ \ell^- \ell^-$, and

\[ X_{5/3} \to H_d^+ u_i, \]

which gives rise to $tbu_i$ final states. We remark that these final states are quite different from the final states usually considered for the $X_{5/3}$ and $t'$ searches \cite{8, 9, 10, 11, 12}. Therefore the heavy quarks of this model are an interesting and motivated incarnation of top partners with non-minimal phenomenology.

Remarkably, we also find that in our model there is a new type of contribution to $\beta\beta_{0\nu}$ decay which not only leads to constraints on the model parameters but also provides a new way to probe our idea.

The paper is organized as follows: in Sec. 2, we present the extension of the supersymmetric type II seesaw model by adding vector-like quark doublets with $Y = \frac{7}{6}$ and show how this makes TeV scale type II natural; in Sec. 3, we discuss phenomenological implications of the model at LHC and in Sec. 4 we point out a new contribution to neutrinoless double beta decay that arises in this model. In Sec. 5 we briefly discuss precision flavor and electroweak observables. In Sec. 6 we give our conclusions.
2 5/3 charge quark and “natural” type II seesaw

2.1 Fine tuning problem in type II seesaw

Type II seesaw [2] mechanism for small neutrino masses is implemented by extending the standard model with the addition of a triplet Higgs field $\Delta$ with $Y = 2$. This leads to the Yukawa coupling in the standard model which has the form:

$$L_Y = h_u \bar{Q}_L H u_R + h_d \bar{Q}_L H d_R + h_{\ell} \bar{L}_L \tilde{H} e_R + f L^T \Delta L + h.c.$$ \hspace{1cm} (1)

The Higgs potential for the model is:

$$V(H, \Delta) = -\mu^2_H H^\dagger H + \lambda_H (H^\dagger H)^2 + M_\Delta^2 \Delta^\dagger \Delta + \lambda_\Delta (\Delta^\dagger \Delta)^2 + \lambda_{H\Delta} (H^\dagger H) (\Delta^\dagger \Delta) + \mu_\Delta H H \Delta + h.c. \hspace{1cm} (2)$$

with $\mu^2_H, M_\Delta^2 > 0$. The simultaneous presence of the couplings $f$ and $\mu_\Delta$ breaks lepton number and upon the EW symmetry breaking the VEV of SM doublet, $\langle H \rangle = v$, triggers the VEV of $\Delta^0$ to be non-zero. We get

$$v_\Delta \equiv \langle \Delta^0 \rangle = \frac{\mu_\Delta v^2}{M_\Delta^2}. \hspace{1cm} (3)$$

As clear from eq. (1) this VEV generates a mass for the neutrinos. There are two ways to get a eV neutrino mass from this formula. Since $\mu_\Delta$ and $M_\Delta$ are both standard model singlet parameters, their values do not depend on the weak scale and can have very high values and it is natural to expect them to be of the same order, i.e. $\mu_\Delta \sim M_\Delta$. Observed neutrino masses require this heavy mass to be of order $10^{14}$ GeV. This keeps the physics of $\Delta$ particles hidden from low energy experiments. On the other hand one could have $M_\Delta \sim$ TeV so that its effects are accessible not only in low-energy high-intensity experiments but also in high-energy colliders. In this case we need to tune $\mu_\Delta \sim 10^{-10}$ GeV or so. It is this fine tuning of $\mu_\Delta$ that we want to address in this paper.

If we consider the supersymmetric version of the triplet seesaw model, the field content of the model is: $Q, u^c, d^c, L, e^c, H_u, H_d, \bar{\Delta}, \Delta$ with a superpotential of the form:

$$W = y_u Q H_u u^c + y_d Q H_d d^c + y_e L H_d e^c + f L^T \Delta L + \mu_H H_u H_d + M_\Delta \Delta^\dagger \Delta + \kappa_u \Delta H_u H_d + \kappa_d \tilde{\Delta} H_u H_u. \hspace{1cm} (4)$$

In this case the triplet VEV for $\Delta$ arises from the potential generated by term $|F_{H_d}|^2$ such that $\mu_\Delta$ of the non-SUSY triplet model becomes $\kappa_u \mu_H$. Since $\mu_H$ is at least of the order of the weak scale to get a realistic EWSB, the fine tuning condition to get small neutrino masses described above translates to the extreme smallness of coupling parameter $\kappa_u \sim 10^{-10}$.

Our goal in this paper is to extend the type II seesaw model so that this fine tuning becomes alleviated.

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1Interestingly, the scalar potential of extensions of the SM with triplets has been shown to significantly impact on the perturbativity and vacuum stability of the Higgs potential. In particular the region of Higgs masses that correspond to a potential stable and perturbative up to very high energy includes lighter masses than what predicted in the pure SM [13].

2Through the mass term $M_\Delta$, when the VEV is generated for $\bar{\Delta}$ also $\Delta$ gets a VEV and vice-versa. Therefore one could get a similar conclusion for the tuning of $\kappa_d$ as well. 
2.2 Vector-like quark extension and alleviation of fine-tuning

We extend the supersymmetric type II seesaw model of the previous section adding a heavy vector-like quark doublet with $Y = 7/6$ so that its component fields, $(X_5/3, t^c)$, have charges $Q(X_5/3) = 5/3$ and $Q(t^c) = 2/3$. We call the two vector-like doublets $Q'$ and $Q'^c$. We also add to the model a gauge singlet field $\sigma$. In order to show that this extension helps to make the TeV scale type II seesaw natural, we impose an additional $Z_4$ symmetry on the model under which $\Delta, \overline{\Delta}, Q'^c, \sigma \rightarrow -\Delta, -\overline{\Delta}, -Q'^c, -\sigma, e^c \rightarrow -ie^c$, and $L \rightarrow iL$ and all other fields are even. In Table 1 we summarize the particle content beyond the MSSM and the charges under the SM gauge group and under the $Z_4$ symmetry. We find the following new Yukawa superpotential:

$$\mathcal{W}' = y_u Q H_u u^c + y_d Q H_d d^c + y_t L H_d e^c + f L^T \Delta L + \mu_H H_u H_d + M_{\Delta} \Delta \Delta + f Q'^c \Delta Q + \lambda Q'^c H_d u^c + y_{Q'} Q'^c Q'^c \sigma.$$

This superpotential is remarkable in that it has interactions of the heavy quarks with several states of the Higgs sector, in particular we remark the interactions with the triplet $\Delta$ and with the Higgs doublet $H_d$. In principle our symmetry allows in the superpotential R-parity and baryon number violating terms $u^c d^c d^c$, which are expected from the fact that our $Z_4$ symmetry is equivalent to lepton number when SM states are considered. In the following we do not consider these type of interactions, assuming for simplicity that they are forbidden by a suitable symmetry. Furthermore we note that from our symmetry follows the absence of the terms $\kappa_u \Delta H_d H_d + \kappa_d \overline{\Delta} H_u H_u$ from the superpotential. As a result, at the tree level, there no induced VEV for the $\Delta$ field. It arises at the one loop level after SUSY breaking from the diagram in Fig. 1. This diagram is induced by

$$F_{Q'} \supset \lambda H_d u^c + y_{Q'} Q'^c \sigma,$$

and generates a finite and calculable coupling $\Delta H_u^* H_d^* \sigma^*$ which gives rise to an effective $\mu_{\Delta, \text{eff}}$
when the singlet field $\sigma$ gets a VEV

$$\mu_{\Delta,\text{eff}} = \frac{y_{Q'}v_\sigma}{16\pi^2} \sum_i y_{ui}f_{qi}\lambda_i \equiv \frac{y_{Q'}v_\sigma}{16\pi^2} Y_{\text{eff}}.$$  \hfill (8)

All squark flavors run into the loop, therefore we define $Y_{\text{eff}} = \sum_i y_{ui}f_{qi}\lambda_i$ which represents the combination of flavored new Yukawa couplings $\lambda_i, f_{qi}$ and SM Yukawa couplings $y_{ui}$ that is relevant for the generation of neutrino masses. For not hierarchical $\lambda_i$ and $f_{qi}$ one can assume that the above formula is dominated by the contribution of the third generation

$$\mu_{\Delta,\text{eff}} \simeq \frac{y_{Q'}v_\sigma}{16\pi^2} y_t f_{q_3}\lambda_3.$$  \hfill (9)

Including this term in the minimization of the potential induces a non-zero VEV for $\Delta^0$ that in general is given by

$$v_\Delta = \mu_{\Delta,\text{eff}}\frac{v_u v_d}{M_\Delta^2},$$  \hfill (10)

which for $Y_{\text{eff}} \simeq y_t\lambda_3 f_{q_3}$ reduces to

$$v_\Delta \simeq \frac{f_{q_3}\lambda_3}{16\pi^2} M_{Q'} \cos \beta \frac{m_t v}{M_\Delta^2}.$$  \hfill (11)

### Table 1: Transformation properties of the field content beyond the MSSM and of the MSSM fields charged under the $Z_4$ symmetry of our model.

| Field | Transformation | $G_{\text{SM}}$ | $Z_4$ |
|-------|----------------|----------------|-------|
| $\Delta$ | $\left( \Delta^{++}, \Delta^+, \Delta^0 \right)$ | (1, 3, 1) | −1 |
| $\bar{\Delta}$ | | (1, 3, −1) | −1 |
| $Q' = \left( X_{5/3} \atop t' \right)$ | (3, 2, $\frac{7}{6}$) | 1 |
| $Q'^{e.c}$ | $(3^*, 2^*, -\frac{7}{6})$ | −1 |
| $\sigma$ | (1, 1, 0) | −1 |
| $L$ | (1, 2, $-\frac{1}{2}$) | $i$ |
| $e^c$ | (1, 1, 1) | $-i$ |


where \( \tan \beta = \langle H_u \rangle / \langle H_d \rangle \), and \( M_{Q'} = y_{Q'} v_\sigma \) is the mass of the exotic quark generated by the VEV of \( \sigma \). For \( M_{Q'} \sim M_\Delta \sim \) TeV at large \( \tan \beta \) and considering only the contribution to neutrino masses from the third generation this corresponds to

\[
v_\Delta \simeq \left( \frac{\lambda_3 f_{q_3}}{10^{-7}} \right) \left( \frac{20}{\tan \beta} \right) \left( \frac{1 \text{ TeV}}{M_\Delta} \right)^2 \left( \frac{M_{Q'}}{1 \text{ TeV}} \right) \text{eV}.
\]

(12)

Let us now recall that current neutrino data for a normal hierarchy imply that the triplet coupling matrix \( f_{ij} \) has the form [14]:

\[
f = \begin{pmatrix}
0.31 - 0.12i & -0.09 + 0.32i & -0.72 + 0.37i \\
-0.09 + 0.32i & 2.53 + 0.04i & 2.19 + 0.01i \\
-0.72 + 0.37i & 2.19 + 0.01i & 3.07
\end{pmatrix}.
\]

(13)

Constraints of the flavor violating process \( \mu \to 3e \) then imply that \( f_{ee} f_{\mu \mu} \leq 4 \times 10^{-5} \) for \( \text{TeV} M_{\Delta^{++}} \). In turn this implies that \( v_\Delta \geq 0.5 \text{ eV} \). Looking at eq. (12), we find that to obtain a realistic mass scale for the neutrino mass generated by the type-II seesaw it is no longer necessary to tune a single coupling to be very small, as it is the case for the generic type-II seesaw model of eq. (5). In fact one can attain the right size of neutrino masses by taking the product of two couplings \( \lambda_3 \cdot f_{q_3} \geq 10^{-7} \). This alleviates significantly the fine tuning problem.

2.3 Quark masses and mixings to \( X_{5/3} \) and \( t' \)

After symmetry breaking, the \( X_{5/3} \) picks up mass of order \( \langle \sigma \rangle \). Guided by the idea that this model describes physics at the weak scale, we choose \( \langle \sigma \rangle \simeq 500 - 1000 \text{ GeV} \). The top partner \( t' \) also picks up the same mass, however it mixes with SM up-type quarks \( u_i \) due to the \( \lambda_i \) and \( f_{q_i} \) couplings and leads to the following mass matrix in the \( u_i - t' \) basis:

\[
M_{u_i - t'} = \begin{pmatrix}
y_{u_i} v_\Delta \\
\lambda_i v_d / \langle \sigma \rangle \\
f_{q_i} v_\Delta / \langle \sigma \rangle
\end{pmatrix}.
\]

(14)

Due to the smallness of the triplet VEV, the mixing of \( t'_L \) to \( u_i \) in general dominates over the mixing of \( t'_R \). The mixing of \( t'_L \) to \( u_i \) is given by \( \theta_{u_i, t'} \equiv \lambda_i v_d / \langle \sigma \rangle \sim 10^{-3} \). We remark that the smallness of this mixing is controlled by the requirement of getting small neutrino masses. Without this restriction, this mixing could be quite significant - although still less than 10-30% or so depending on what the mass of the \( X_{5/3} \) quark is. From this mass matrix, we also learn that the splitting between \( X_{5/3} \) and \( t' \) is very small (less than a GeV). The rest of the quarks and leptons get their masses as in the standard model. The only effect on the standard model quark sector is that there will very small mixings between each quark and the \( t' \).

Besides introducing a mixing between new matter and SM matter, the coupling \( \lambda \) has its most apparent effect in the modification of the decay of the heavy quarks with respect to the minimal models of heavy quarks, where they usually decay to Goldstone bosons in the form of longitudinal SM gauge bosons. As remarked above, in principle also \( f_q \) mediates a mixing, but it is significantly suppressed by the smallness of the VEV \( v_\Delta \). We will see below that
the interactions mediated by \( f_q \), as well as the one mediated by \( \lambda \), have several interesting phenomenological implications for the decay of the new fields of our model. In particular they induce new decay modes for the heavy quarks, making them phenomenologically distinct from the heavy quarks of minimal models.

3 Decays of the \( \frac{5}{3} \) charge quark and the seesaw triplet

The states \( \Delta \) and \( Q' \) have a peculiar phenomenology, especially the \( X_{5/3} \) and the \( \Delta^{++} \) states. In fact their exotic electric charge gives rise to specific signatures characterized for instance by the presence of leptons of same sign that can also have a resonant structure. These features are indeed exploited in theoretical and experimental studies on the observation of these states [7, 3, 15, 8, 9, 10, 11].

Despite the large literature about the states \( \Delta \) and \( Q' \) the phenomenology of models where both these states appear is less studied. We have seen above how the simultaneous presence of these states can induce neutrino masses radiatively. For this goal are necessary several interactions beyond those of the SM and the MSSM. In particular from eq. (12) we see that three couplings of the model of eq. (5) are involved in the neutrino mass generation:

- \( \lambda \) which mediates interactions among \( Q', u^c \) and \( H_d \) and that after EWSB it also causes the \( t' - u_t \) mixing;
- \( f \) which gives interactions among lepton doublets and the Higgs triplet \( \Delta \);
- \( f_q \) which gives interactions among \( Q', Q \) and the Higgs triplet \( \Delta \).

Only the product of these couplings is constrained by neutrino masses, \( \lambda f f_q \simeq 10^{-7} \), so that there is still significant freedom for the value of each single coupling. The masses of the new states \( Q' \) and \( \Delta \) are expected to be of the same order and there is no substantial preference for one being heavier than the other. We first consider the case \( M_{X_{5/3}} > M_{t'} > M_{\Delta} \), which is somewhat more interesting than the converse because in this case scalar non-colored states can appear in the decay chain of copiously produced colored objects. In this case the exotic \( X_{5/3} \) has several possible decay modes discussed in the following.

Weak gauge interactions mediate decays into states of the \( Q' \) multiplet in the same way they do in typical models where heavy quarks appear:

\[
X_{5/3} \rightarrow W^+ t',
\]
\[
X_{5/3} \rightarrow W^+ t,
\]

where the latter involves a \( t - t' \) mixing. The first process is kinematically forbidden for most of the parameter space of our model or is generically suppressed for virtual \( W \) by the small mass difference between \( X_{5/3} \) and \( t' \) or by the suppression in the \( t - t' \) mixing from the small \( v_d/\langle\sigma\rangle \) that we discussed in Section 2.3.

\[3\text{In fact one can even find discrete symmetries slightly different from our } Z_4 \text{ that would give both the mass of } \Delta \text{ and that of } Q' \text{ from renormalizable Yukawa superpotential interactions. In that case the resulting mass for the two states is controlled by two couplings with no a priori hierarchy.}\]
Our model is characterized by the existence of decay modes mediated by $\lambda$

$$t' \rightarrow H^0_d u_i, \quad (17)$$

$$X_{5/3} \rightarrow H^+ d u_i, \quad (18)$$

where $H^+_d (H^0_d)$ is a component of the charged (neutral) Higgs boson that is mostly made out of the doublets 4.

Interestingly, these modes can be the dominant ones if one takes $\lambda$ to be the largest coupling and ascribes the smallness of neutrino masses to $f \cdot f_q$. In fact we estimate that the coupling $\lambda$ could be of order 0.1 without causing tensions in precision observables in the $Z \rightarrow b \bar{b}$, $Z \rightarrow$ hadrons decays, and $D$ meson mixing [16], which are potentially sensitive to this coupling 5. Depending on the way the states $H^+_d$ decay, these decay modes could be different enough from the usual decay modes assumed in searches for heavy quarks and they may even cause a significant relaxation of the bounds on the existence of $Q'$.

In the triplet-doublets decoupling limit, and for moderate or large $\tan \beta = (H^0_u)/(H^0_d)$, the field $H^0_d$ is very close to the heavy mass eigenstates of the MSSM scalar spectrum $H, A$. Therefore the dominant $t'$ decays mediated by $\lambda$ are expected to be

$$t' \rightarrow H u_i, \quad (19)$$

$$t' \rightarrow A u_i, \quad (20)$$

whereas $t' \rightarrow h u_i$ is likely to be sub-dominant due to the value of $\tan \beta$. In principle the decay rate may depend on the flavor of the up-type quark involved in eq.(19) and (20). We expect

$$H, A \rightarrow b \bar{b}$$

to be among the dominant decay modes, and likely the dominant decay mode, of the heavy Higgs bosons. In fact decays into Higgs bosons and gauge vectors are typically suppressed by phase-space or $\tan \beta$. The decay in fermionic partners of the electroweak gauge bosons

$$H \rightarrow \chi^0 \chi^0, \chi^+ \chi^-$$

are in principle possible, although the decay rate is highly dependent on mass and mixing parameters of the neutralino-chargino sector. In any case it is worth stressing the possibility. In fact a significant branching fraction in this decay mode would result in a significant change of the observable signature of the production of the heavy quarks.

Sticking to the limit where the triplet is decoupled from the doublets, the $X_{5/3}$ decays

$$X_{5/3} \rightarrow H^+ u_i \rightarrow t \bar{b} u_i \quad (21)$$

4In this model the VEV of the triplet field is small and the trilinear terms of the type $\Delta H H$ are forbidden at tree level, therefore it is rather natural to assume that the triplets and the doublets do not mix significantly.

5Ultimately these effects are small thanks to the smallness of the $u_i - t'$ mixing described in Section 2.3 and the mass of the $Q'$ fields in the TeV range. It should also be noted that this theory predominantly generates chirality conserving operators for the FCNC processes. See a more extended discussion in Sec. 5.
where the decay of the charged Higgs boson into fermions is favored by the smallness of $v_d$ and the decay into $tb$ is assumed to be kinematically allowed $^6$. The hierarchy between $v_u$ and $v_d$ suppresses also other decays such as $X_{5/3} \rightarrow W^+ u_i$, making eq. (21) the main decay mode of $X_{5/3}$ when $\lambda$ dominates.

We see that having two Higgs doublets plays an essential role in the phenomenology of the exotic fermions that can now decay in the extra physical degrees of freedom of the scalar sector (see Refs. [19, 20, 21] for a recent related discussion about heavy quarks phenomenology in the context of models of strongly interacting EWSB). The different decay pattern of these heavy quarks is likely to make the standard searches less sensitive to the type of object present in this model. However, we do not expect this scenario to be free from experimental bounds. In fact searches for resonances that decay in $tb$ final state [22, 23, 24] are potentially sensitive to the production of $H^+$ in cascades of the exotic quarks as in the decay eq. (21). Furthermore the searches for heavy Higgs bosons of the MSSM [25, 26] are in principle sensitive to the decays of heavy quarks into states of the extended Higgs sector. Also, an extension of searches for top quarks in association with Higgs bosons of some more general scope of the current ones [27] would certainly allow to investigate in full generality the phenomenology of heavy quarks.

The coupling $f_q$ mediates several decays of $Q'$ that are distinctive of our model. We estimate that a coupling of order order 0.1 would respect available constraints from data meson oscillation [16] and $Z$ pole precision observables [38]. Therefore we can consider a scenario where the decay of $Q'$ is dominated by the coupling $f_q$ while the smallness of the neutrino mass is ascribed to the couplings $\lambda$ and $f$. The coupling $f_q$ mediates

$$X_{5/3} \rightarrow \Delta^{++} d_i ,$$

$$X_{5/3} \rightarrow \Delta^+ u_i ,$$

with strength that in principle may depends on the flavor of the down-type and up-type quarks $d_i$ and $u_i$. The mode of eq. (22) is new and distinctive of the present model with respect to the usual models where the state $Q'$ is not involved in neutrino mass generation and therefore it does not have the interaction $f_q Q' \Delta Q$. The way this decay can be detected in practice depends a lot on the decay of the $\Delta$, which we discuss below. The mode eq. (23) and other decay modes of $t'$ mediated by $f_q$ such as

$$t' \rightarrow \Delta^+ d_i ,$$

$$t' \rightarrow \Delta^0 u_i ,$$

are new as well. However, as they involve Higgs states of charge one or zero, they are likely to result in final states quite similar to those of $t' \rightarrow Z t$, $ht$ and $X_{5/3} \rightarrow W^+ t$ and so they are less distinctive of this model. In case of observation of an excess in the searches sensitive to these decay modes one should take into account the possibility of extra contributions, with possibly peculiar flavor structure, that contribute to the lepton rich final states of these BSM searches [8, 9, 10, 11]. Due to the similarities among the standard decay modes considered

$^6$Although it is not strictly forbidden to have lighter charged Higgs bosons, we focus on the case of a charged Higgs heavier than 350 GeV such to straightforwardly satisfy the limits on the $b \rightarrow s\gamma$ decay rate [17] and $t \rightarrow H^+ b$ searches [18].
| Decay                        | Rate                      | Dominant for               |
|------------------------------|---------------------------|---------------------------|
| $M_{Q'} > M_\Delta$          |                           |                           |
| $Q'_\chi H_d u^c$            | unsuppressed              | $\lambda \gg f_q$         |
| $Q'_\tau \Delta q \rightarrow LL$ | unsuppressed              | $f_q \gg \lambda$         |
| $Q'_\tau V u^c$              | suppressed by $\frac{\lambda v}{\langle \sigma \rangle}$ | $\lambda \gg f_q$ and $m_{H_d} > m_{Q'}$ |
| $M_\Delta > M_{Q'}$          |                           |                           |
| $\Delta \rightarrow LL$     | unsuppressed              | $f \gg \lambda, f_q$      |
| $Q'_\chi H_d u^c$            | unsuppressed              | $f_q, \lambda \gg f$      |
| $Q'_\tau q_L \Delta^* q_L LL$ | 3-body suppressed         | $f_q \gg \lambda, f$      |
| $Q'_\tau V u^c$              | suppressed by $\frac{\lambda v}{\langle \sigma \rangle}$ | $f_q \gg f$ and $m_{H_d} > m_{Q'}$ |

Table 2: Summary of the decay patterns of the non-SUSY particles of our model. The shorthands $L = \ell, \nu$, and $V = Z, W$ must be spelled out in the suitable combination for each final state. For the decay of $H_d$ in the large tan $\beta$ limit and neglecting the small doublet-triplet mixing the decays $H_d \rightarrow f \bar{f}$ or $\chi \bar{\chi}$ are both possible, but the $f \bar{f}$ final state is generically favored.

in the searches of heavy quarks and the decays eq. (24) we can take the current bounds from [8, 9, 10, 11] and future updates as a good estimate of the mass bound on $m_{Q'}$ for the scenario where $f_q$ dominates the decay of $Q'$.

The triplet $\Delta$ decays are quite well studied in the context of type II seesaw model, however the picture that emerges in this model can be substantially different. In a way similar to the standard type II seesaw model the gauge interactions mediate

$$\Delta^- \rightarrow W^- W^-,$$
$$\Delta^- \rightarrow W^- \Delta^-,$$

that are suppressed by the VEV $\langle \Delta^0 \rangle$ and by the smallness of the typical splitting of the components of the $\Delta$ triplet. Therefore these interactions do not typically have impact of the phenomenology of $\Delta$ and we neglect them.

It should be remarked that being a SUSY model there are two Higgs doublets, and therefore more physical degrees of freedom into which the $\Delta$ can decay. In fact, in addition to the listed modes there could be interactions with the additional scalars $H^\pm, H, A$ of the MSSM. For the decay the most important interactions should arise from trilinear terms. These are either originating from quartic term of the form $\Delta^2 H^2$ or from genuine cubic terms $\Delta H H$. In both cases these interactions are suppressed. In fact the
As in the standard type II seesaw model, the interaction $f$ mediates the decays into leptons that are used in the current searches for the seesaw triplet [28]

\[
\begin{align*}
\Delta^{--} & \rightarrow \ell\ell, \\
\Delta^- & \rightarrow \ell\bar{\nu}, \\
\Delta^0 & \rightarrow \nu\bar{\nu}.
\end{align*}
\]

When the decay of $\Delta$ is dominated by $f$ the searches for doubly charged resonances [28] are in force. Depending on the lepton flavor combination that appear in eq. (27) and eq. (28) a bound on the productions cross-section varying from about 1 fb to about 100 fb is in force. Considering that the source of the doubly charged scalar in this model would be the colored fermion $X_{5/3}$ this translates into a bound on the mass of the $Q'$ up to 1 TeV. Therefore in this model the bounds from the peculiar decay mode eq. (22) are stronger than the bounds on heavy quark from searches that look for the signatures of minimal models of heavy quarks [8, 9, 10, 11].

Differently from the standard type II seesaw model without the $Q'$ fields, in our model $\Delta$ can decay via the interactions mediated by the couplings $\lambda$ and $f_q$. For instance the doubly charged scalar has the decay mode

\[
\Delta^{++} \rightarrow \bar{d}_i X^*_{5/3} \rightarrow \bar{d}_i u_j H^+ \rightarrow \bar{d}_i u_j t \bar{b}
\]

which involves the product of couplings $\lambda f_q$. This mode can dominate when the coupling $f$ is by far the smallest among $\lambda, f_q, f$. A very small $f$ may be suggestive of some tuning of the couplings of the model, going against the original motivation for our model. However the resulting phenomenology of eq. (30), is quite different from the case where $f$ dominates the decay of $\Delta$, and therefore we believe it deserves some attention. Furthermore, especially for light $\Delta$, the coupling $f$ may be required to be small to comply with lepton number violation bounds from $\mu \rightarrow eee$ and radiative processes such as $\mu \rightarrow e\gamma$ [29].

For the case where the triplet is heavier than the exotic quark, $M_\Delta > M_{Q'}$, there are significant differences. Most importantly the production of the electroweak states of $\Delta$ cannot be enhanced by the production of $\Delta$ in decay chains of colored states. However one can still have the doublet-like states $H_{d,u}$ in such decay chains, which is still a interesting signature of this type of heavy quarks.

In the case $f \gg \lambda, f_q$ the triplet $\Delta$ decays as in the usual type-II seesaw model. In particular there is a decay

\[
\Delta^{--} \rightarrow \ell\ell
\]

and the bounds from [28] are in force.

In the case $f \ll \lambda, f_q$ the triplet cannot decay into leptons and it has to decay through the coupling $f_q$ into a left-handed SM quark and a heavy quark $Q'$

\[
\Delta \rightarrow Q' q_{L,SM}.
\]

quartic terms affect the decay of $\Delta$ only after one of the $\Delta^0$ has taken a VEV, which we know must be a small quantity in this model. On the other hand trilinear terms of the form $\Delta HH$ are forbidden at tree-level in this model, hence do not in general change the phenomenology of the decay of $\Delta$. 

12
Therefore the decay of $\Delta$ is effectively giving rise to an extra source of $Q'$. This extra source of $Q'$ is certainly sub-dominant to the QCD production of $Q'$. However we can exploit this fact to use the phenomenology of $Q'$ as a single probe to investigate both the existence of $Q'$ and $\Delta$.

After the production, either through QCD gauge interactions or in the decay of $\Delta$, the decay of $Q'$ can either be ruled by $\lambda$, $f_q$ or gauge interactions. The decay through the coupling $\lambda$ is the only one that can go through without suppression. In fact it gives a two-bodies decay $Q' \to H_d u^c$ that results in signals similar to eqs. (19),(20),(21) discussed above. These decays, when kinematically allowed, are expected to be dominant because the decays mediated by $f_q$ or gauge interactions pay either a phase-space or a mixing suppression. As stressed already in the discussion of the case $M_{Q'} > M_\Delta$ the decay modes eqs. (19),(20),(21) give rise to interesting signatures that characterize our type of non-minimal heavy quarks. Bounds on these signatures can be obtained recasting (or extending) current searches as detailed above.

The decay mediated by $f_q$ gives a 3-bodies decay into an off-shell $\Delta$. The virtual $\Delta$ can decay only into leptons, that are the only light final state available and overall we have a decay of $Q'$

$$Q' \to \Delta^* q_{L,SM} \to \ell\ell q_{L,SM}. \quad (32)$$

This signature is in general very similar to the one with non-resonant same sign and opposite sign leptons investigated in the searches for minimal heavy quarks [8, 9, 10, 11]. Therefore we expect that, up to differences in the efficiencies, these searches would constrain the scenarios where $Q'$ decays in three-bodies via the coupling $f_q$. This decay mode is expected to dominate when either the states $H, A, H^+$ are too heavy for the decays eqs. (19),(20),(21) to be on-shell, or there is a large hierarchy of couplings $f_q > 16\pi^2\lambda$.

The 3-bodies decay mediated by $f_q$ has to compete with the gauge decay modes, which are the same as in eq. (16). In this case there is a suppression of order $v_d/\langle\sigma\rangle$ from the $u_i - t'$ mixing angle, which is of the same size of the 3-body phase space suppression. Therefore the dominance of the gauge modes with respect to the 3-bodies decays of eq. (32) depends on the detailed hierarchy among the couplings of the model. In any case, whether are the gauge modes or the 3-bodies modes of eq. (32) that dominates the decay, the standard searches [8, 9, 10, 11] should apply up to minor changes in the mass reach.

A summary of the decay of non-SUSY states is presented in Table 2.

### 3.1 SUSY modes

Given the presence of SUSY partners we consider the phenomenology of these particles in our model, focusing mostly on the new features with respect to the MSSM and SUSY seesaw model. We have a SUSY partner of the seesaw triplet, that we denote as tripletino $\tilde{\Delta}$, and scalar partners of the heavy quarks, the scalar heavy quarks $\tilde{Q}^'$ and $\tilde{Q}^{'+c}$.

When the coupling $f$ dominates among the new couplings, our model is essentially reproducing the phenomenology of the SUSY seesaw models that have been studied for instance in [30, 31, 32]. The typical signature of the production of $\tilde{\Delta}$ has several leptons and missing transverse energy, analogously to the final states considered in [33, 34]. Similar final state are sensitive to the gauge decays of the tripletino as for instance $\Delta^{-}\to W^{-}\Delta^{-}$. It is also
Decay Rate Dominant for

\[ M_{Q'} > M_\Delta \]

- \( \tilde{Q}_{\chi}' \tilde{H}_d u^c \) unsuppressed \( \lambda \gg f_q \)
- \( \tilde{Q}_{\tilde{\chi}}' \tilde{H}_d \tilde{u}^c \) unsuppressed \( \lambda \gg f_q \)
- \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET unsuppressed \( f_q \gg \lambda \)
- \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET unsuppressed \( f_q \gg \lambda \)
- \( \tilde{Q}_{\tilde{g}}' \tilde{V} \tilde{u}^c \) and \( \tilde{Q}_{\tilde{g}}' \tilde{\lambda}_i u^c \) suppressed by \( \frac{\lambda v_d}{\langle \sigma \rangle} \) \( \lambda \gg f_q \)

\[ M_\Delta > M_{Q'} \]

\[ \tilde{\Delta}_{\tilde{f}_q} \tilde{Q}_{\tilde{f}_q}' \]

- \( \tilde{Q}_{\chi}' \tilde{H}_d u^c \) unsuppressed \( f_q, \lambda \gg f \)
- \( \tilde{Q}_{\tilde{\chi}}' \tilde{H}_d \tilde{u}^c \) unsuppressed \( f_q, \lambda \gg f \)
- \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET 3-body suppressed \( f_q \gg \lambda, f \)
- \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET 3-body suppressed \( f_q \gg \lambda, f \)
- \( \tilde{Q}_{\tilde{g}}' \tilde{V} \tilde{u}^c \) suppressed by \( \frac{\lambda v_d}{\langle \sigma \rangle} \) \( \lambda \gg f_q \gg f \)
- \( \tilde{Q}_{\tilde{g}}' \tilde{\lambda}_i u^c \) suppressed by \( \frac{\lambda v_d}{\langle \sigma \rangle} \) \( \lambda \gg f_q \gg f \)

\[ \tilde{\Delta}_{\tilde{f}_g} \tilde{L}, \tilde{W} \tilde{\Delta}' \]

leptons + mET unsuppressed \( f \gg \lambda, f_q \)

| Decay Pattern | Rate | Dominant for |
|---------------|------|--------------|
| \( \tilde{Q}_{\chi}' \tilde{H}_d u^c \) | unsuppressed | \( \lambda \gg f_q \) |
| \( \tilde{Q}_{\tilde{\chi}}' \tilde{H}_d \tilde{u}^c \) | unsuppressed | \( \lambda \gg f_q \) |
| \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET | unsuppressed | \( f_q \gg \lambda \) |
| \( \tilde{Q}_{\tilde{f}_q}' \tilde{q}_L \tilde{\Delta} \rightarrow \) leptons + jets + mET | unsuppressed | \( f_q \gg \lambda \) |
| \( \tilde{Q}_{\tilde{g}}' \tilde{V} \tilde{u}^c \) and \( \tilde{Q}_{\tilde{g}}' \tilde{\lambda}_i u^c \) | suppressed by \( \frac{\lambda v_d}{\langle \sigma \rangle} \) | \( \lambda \gg f_q \) |

Table 3: Summary of the decay patterns of the SUSY particles of our model. The shorthands \( L = \ell, \nu \), and \( V = Z, W \) must be spelled out in the suitable combination for each final state. For the decay of \( H_d \) in the large \( \tan \beta \) limit and neglecting the small doublet-triplet mixing the decays \( H_d \rightarrow \bar{f}f \) and \( \chi \bar{\chi} \) are both possible, however the decay into \( \bar{f}f \) is generically favored. The decays of the SUSY states higgsino \( \tilde{H}_d \) and gaugino \( \tilde{\lambda}_i \) are more model dependent and we leave them unspecified here.

\(^{a}\) For this mode to dominate \( m_{H_d} > m_{Q'} \) is required, otherwise the \( H_d^0 \) modes are kinematically allowed and unsuppressed.
possible to have resonant pair production of leptons mediated by gauge interactions in the
tripletino SUSY decay $\Delta^{++} \rightarrow \Delta^{++}\chi^0$ for which standard [28] and dedicated searches [35]
apply. We expect that with minor modification due to the efficiencies the results of the LHC
searches can be applied for this models as well.

The true nature of our model is best displayed when the couplings that dominate the
declay of $\Delta$ and $\tilde{Q}$ are either $f_q$ or $\lambda$. As in the case of the decay of non-SUSY particles
described above, when the triplet $\Delta$ does not decay to leptons it produces heavy quarks (or
SUSY partners of them) so it is best to start the discussion for $M_\Delta > M_{Q'}$. In this case the
declay
$$\Delta \rightarrow q\tilde{Q}'$$  \hspace{1cm} (33)
is dominant mode when $f_q \gg f$. In fact the other SUSY mode $\tilde{\Delta} \rightarrow \tilde{q}Q'$ is likely to be
suppressed, or even inaccessible, because of kinematics, as both $\tilde{q}$ and $Q'$ are heavy states.
At this point the resulting signature is determined by the decay of $Q'$. For $f_q \gg \lambda$ the
hierarchy of couplings can win over the price of emitting an off-shell $\Delta$ and the resulting
declay is
$$\tilde{Q}' \rightarrow \tilde{q}\Delta^{*} \rightarrow \text{mET + jets + leptons},$$  \hspace{1cm} (34)
where the off-shell nature of the $\Delta$ implies non-resonating leptons in the final state. A similar
final state is obtained when $Q'$ decays in the SUSY partner final state $\tilde{Q}' \rightarrow q\tilde{\Delta}^{*}$.

When the coupling $\lambda$ dominates the decay of $\tilde{Q}'$ we have
$$\tilde{Q}' \rightarrow \tilde{u}cH_d \rightarrow \text{mET + jets + H}_d,$$  \hspace{1cm} (35)
where, depending on the charge of $Q'$, the resulting state in $H_d$ decays to $bb$ or $tb$. The decay
in the higgsino final state is less well defined because the the SUSY partner $H_d$ decays are
more model dependent. Therefore we leave them unspecified here, although we expect them
to be similar to those included in a inclusive search of for the final states of eq. (35).

In principle the decay of $Q'$ can result in MSSM squarks and gauginos production via the
declay $\tilde{Q}' \rightarrow \tilde{q}V$ and $\tilde{Q}' \rightarrow q\chi$, however these decays are suppressed by the mixing between
exotics and SM quarks and therefore they usually have little importance.

The discussion is very similar for the case $M_\Delta < M_{Q'}$, that essentially differs from the case
$M_\Delta < M_{Q'}$ for the effect of off-shell phase-space factors. This has important implications
for the case of the chain eq. (34) where the presence of resonating leptons can make a big
difference for the experimental searches.

A summary of the decay of SUSY states is presented in Table 3.

A comment on the bounds from standard SUSY searches in our model is in order. In fact
we have new sources of SUSY particles and new interactions that in principle can change
the bounds for our model. Generically we expect the current SUSY searches to give more
stringent bounds in our model on the existence of super-partners. In fact in our model
there are extra sources of squarks such as for instance the decays $\tilde{Q}' \rightarrow \tilde{q}\Delta$ and $\tilde{\Delta} \rightarrow \tilde{q}Q'$.
In most cases these decays have small branching fractions, because the same particles can
decay without suppression in other modes that have at least one light final state, hence they
systematically beat the modes containing a squark that have always two heavy modes. For
In this reason we expect only a modest overproduction of quarks and therefore limits are only mildly affected by the new dynamics of our model.

In principle the decay of MSSM states receive corrections from the existence of new states and new interactions that characterize our model. A particularly relevant case is that of the squarks that can decay $\tilde{q} \rightarrow \Delta \tilde{Q}'$ and $\tilde{q} \rightarrow \tilde{\Delta} \tilde{Q}'$, that in general gives rise to more lepton-rich final states than the usual squark decays of the MSSM. From this we expect that a slightly harder bounds from the LHC searches could apply to our model. However, a precise determination of the new bound requires a detailed study. In fact there are other effects that contribute against the isolation of the signal with the strategies adopted in current searches. For instance our final states typically result in a larger multiplicities and this is not always beneficial for the identification of new physics, as in general the individual objects in high multiplicity events tend to be softer.

4 New contributions to neutrinoless double beta decay from 5/3 charge quarks

Neutrinoless double beta decay [36] is a powerful probe of new physics that gives rise to violations of lepton number. Although the experiments probe relatively low energies in the domain of nuclear physics, the intensity and the precision of these experiments is such as to probe physics at much higher scales. For instance ascribing the $0\nu\beta\beta$ transition to a dimension-9 operator the current generation of experiments is sensitive to operators suppressed by roughly $(1/ \text{TeV})^5$. This means that neutrinoless double beta decay can also probe models of EWSB when the breaking of the gauge symmetry of the SM is in some manner entangled with the breaking of lepton number. Electroweak models for the generation of neutrino masses are therefore an excellent target for $0\nu\beta\beta$ experiments. Our model is no exception to this statement, as we have demonstrated that, not only the neutrino mass is ultimately generated by the VEV of Higgs particles, but also we have found that the overall
scale of neutrino masses is set by the scale of SUSY breaking, hence by the electroweak scale.

In our model there is a new short-range contribution to neutrino-less double beta decay that involves the couplings f_q, f, and λ via the diagram in Fig. 2. Remarkably it involves the simultaneous propagation of an exotic quark and a doubly charged scalar. This diagram gives rise to 0νββ transition mediated by an effective operator of dimension 9, whose strength is partly correlated to the overall scale of neutrino mass. To the best of our knowledge this short-range contributions has not been explicitly spelled out before, although the resulting dimension-9 operator that mediates double beta decay and a possible set of mediators has appeared in studies of the possible decompositions of the 0νββ operator [37]. Therefore this type of contribution appears to be characteristic of seesaw models that have heavy exotic quarks in their spectrum. The presence of this new characteristic interplay between particles in the reach of collider and low energy processes definitively constitute further motivation to study our model.

We estimate that the new contribution to neutrinoless double beta decay mediated by X_{5/3} results in a matrix element

\[ M_{X_{5/3}}^{\beta\beta} \simeq \frac{G_F f_{q_1} \lambda_1 f_{ee}}{2 M_\Delta^2 M_{Q'} M_{Q'}} \bar{e}_L C e_L \bar{u}(1 + \gamma_5) \gamma_\alpha d \bar{u} \gamma_\alpha (1 + \gamma_5) d, \]  

where we treat collectively the mass of t' and X_{5/3} by denoting them as M_{Q'}. We remark that the rate for double beta neutrinoless decay depends on combination of the new couplings f, λ, and f_q that is in general different from the combination Y_{\text{eff}} that determines the mass scale of the neutrino. This is remarkable in that it allows to disentangle the rate of the neutrinoless double beta decay from the absolute scale of neutrino masses. Using eqs. (8)-(12) the coefficient of this matrix element can be written in terms of the neutrino mass so that

\[ M_{X_{5/3}}^{\beta\beta} \simeq G_F 8\pi^2 \frac{m_\nu}{m_t M_{Q'}^3} \frac{\lambda_1 f_{q_1}}{\lambda_3 f_{q_3}} \bar{e}_L C e_L \bar{u}(1 + \gamma_5) \gamma_\alpha d \bar{u} \gamma_\alpha (1 + \gamma_5) d, \]  

where we assumed the coupling f to not be hierarchical in flavor space, as suggested by the large neutrino mixing, and we also assumed the sum in eq.(8) to be dominated by the third generation, i.e. Y_{\text{eff}} \simeq \lambda_3 f_{q_3} y_t \geq \lambda_i f_{q_i} y_{u_i} \text{ for } i = 1, 2. This gives a coefficient that scales like

\[ 10^{-15} G_F^2 \frac{\lambda_1 f_{q_1}}{\lambda_3 f_{q_3}} \left( \frac{m_\nu}{0.1 \text{ eV}} \right) \left( \frac{1 \text{ TeV}}{M_{Q'}} \right)^3 \text{ GeV}^{-1}. \]  

Note that the smallness of \( \mu_\Delta \) in eq.(12) allows for \( \lambda_3 f_{q_3} \) to be small and still to dominate the generation of neutrino masses. Considering only the first and third generation involved here one finds that the third generation dominates in the neutrino mass generation as long as

\[ \frac{\lambda_3 f_{q_3}}{\lambda_1 f_{q_1}} > \frac{y_t}{y_u} \simeq 10^{-5}. \]  

Current experimental bounds on the life time for this process roughly translate into bound on the coefficient of this operator to be less than about \( G_F^2 \times 10^{-8} \text{ GeV}^{-1} \) [36, 38]. Thus there are regions of the parameter space of the model where the double beta amplitude can be
near the current upper limit. For instance, if one assumes safe couplings for precision flavor and electroweak observables, say $\lambda_{1,2} \sim 0.1$ and $f_{q_1, q_2} \sim 0.1$, taking an “inverse” hierarchy for the couplings $\lambda_i$ and $f_{q_i}$, say $f_{q_3} \sim \lambda_3 \sim 10^{-4}$ or so, the double beta amplitude for TeV heavy quarks is only about one order of magnitude below the current limit and at the same time eq. (12) predicts sub-eV masses for the neutrinos. Therefore it is possible to test this model through future searches for neutrinoless double beta decay. We stress that such hierarchy of the couplings $f_{q_i}$ and $\lambda_i$ can be cross-checked because it is observable in the lack of heavy flavors in the decay of the heavy quarks summarized in the Tables 2 and 3.

Furthermore we find that the signal in neutrinoless double beta decay is correlated in a peculiar way to signals in deep inelastic scattering (DIS) experiments. In fact at momentum scale of order $m_W$ the effective operator that generates $0\nu\beta\beta$ in our model opens up leaving a contact interaction $udWee$ from the interactions of the heavy degrees of freedom $X_{5/3}$ and $\Delta^{++}$. This interaction mediates transitions

$$e^- u \rightarrow d^+ W^- ,$$

which gives rise to a hard lepton of charge opposite to that of the beam in the DIS experiment. This process is very interesting. In fact it can in principle be used to test the flavor structure of the $f$ coupling responsible for the seesaw. Furthermore, the rate of this reaction (especially for the positron flavor in the final state) is tightly correlated to the rate of $0\nu\beta\beta$. Therefore such DIS experiment allows to test the fraction of the $0\nu\beta\beta$ rate that is due long-range contributions such as the neutrino mass versus the fraction that comes from short-range contributions such as the one of our model in Figure 2. Similar correlations exists in model of leptonic number violating R-parity violating SUSY [39, 40], where, however, sleptons are expected in place of the $W$ boson of our process eq. (40). Discerning $W$ bosons from sleptons in DIS final states, of course, offers further information on the actual origin of a possible $0\nu\beta\beta$ signal observed in the next round of experiments.

### 5 Other implications of the model

The presence of new quarks that mix with the SM quarks gives rise to several changes in the phenomenology at energies even much below the mass the of the new quarks.

For instance, the couplings $f_{q_i}$ and $\lambda_i$ have a flavor index and therefore these couplings can lead not only to departure from CKM unitarity (which are less than a percent level), but they can also lead to flavor changing neutral currents (FCNC) e.g. $K - \bar{K}$ oscillations mediated by box graphs involving the exchange of $\Delta$, $H_d$ and $Q'$. However, this contribution to the $\Delta S = 2$ effective Hamiltonian comes with a strength $\sim \frac{f_{q_1}^2 f_{q_2}^2}{16\pi^2 M_Q} \simeq 10^{-6}$ TeV$^{-2}$ for $f_q \sim 0.1$ and $M_Q \sim 1$ TeV, which is right at the current bound for chirality and CP conserving transitions [16]. Similar contributions also arise for $B$ and $D$ mesons mixing and are more easily compatible with the bounds. Along the same lines one can find analogous conclusions for to the case where only the coupling $\lambda$ is used in the box diagram. Since chirality flipping operators have in general much more stringent bounds [16] it is worth remarking that in our model it is likely that this kind of operators are generated with an extra suppression with respect to the chirality conserving ones. In fact these chirality flipping operators can be
generated only using simultaneously both the couplings $\lambda_i$ and $f_{q_i}$. As displayed in eq. (12) we need the product of two couplings to be $\lambda_3 f_{q_3} \sim 10^{-7}$ to reproduce the overall scale of neutrino masses. Therefore it is natural to assume that at least one of the two couplings must be small, such that the main effect in FCNC comes from box diagrams where only the large coupling appears. Alternatively, chirality flipping operators arise from mixed box diagrams with one $W$ boson and one new scalar, e.g. $\Delta^+ W$ box graphs. These are in general suppressed as they involve small VEVs.

Furthermore, new quarks with same electric charge and different $SU(2)$ quantum numbers than the ordinary quarks modify the couplings of the SM quarks to the weak gauge bosons. Especially the couplings of the $Z$ boson are very accurately measured and there are constraints on the amount of mixing that can be tolerated. As discussed in Section 2.3, in our model the mixing between SM and new quarks is in general small, due to the smallness of the VEVs responsible for the mixing, that are $v_d$ and $v_\Delta$ for the $t_R$ and $t_L$ mixing, respectively. The new quarks also contribute to oblique parameters $S$ and $T$ via their mixing with the SM quarks, which provide yet another constraint. Recently these issues have been analyzed in the light of the most up to date results on precision observables and direct searches of heavy quarks [41] finding that for small mixings as in our model the prediction of the electroweak observables are generically in agreement with the experiment.

From the above discussion it is clear that precision flavor and electroweak observable do not deviate significantly from the SM prediction once the new particles have TeV mass and the new couplings $\lambda$ and $f_q$ are of the order 0.1 or less. Incidentally we remark that, unless the model is complicated such to have cancellations in the BSM contribution to flavor and electroweak precision observables [42], this implies that the single production at colliders of the new states through the couplings $\lambda$, $f_q$ is generically much suppressed with respect to the pair production via gauge interactions.

6 Conclusions

We have presented a model of physics of the electroweak scale where the weak scale is originated from the breaking of supersymmetry. At the same time in our model the breaking of supersymmetry triggers the generation of neutrino masses through a mechanism much like the seesaw of type-II. Our model is distinct from the usual SUSY seesaw in that the size of the trilinear interaction of the seesaw triplet and the Higgs doublets, a term $\Delta H_u^* H_u$ in our case, is not put by hand to a tuned value, but rather it is dynamically connected to the scale of the breaking of SUSY, hence to the electroweak scale. The connection between the scale of generation of neutrino mass and that of stabilization of the Higgs potential is embodied by the diagram of Figure 1. To establish this link between the weak scale and neutrino mass generation is necessary to extend the SUSY type-II seesaw model with a new quark of hypercharge 7/6, whose interactions, upon SUSY breaking, make the neutrino mass non-vanishing.

Very remarkably, our mechanism for the generation of the $\Delta HH$ term of the type-II seesaw model has considerably less fine-tuning that the standard mechanism. As shown in eq. (12), in our model the smallness of this trilinear term, and therefore of the neutrino mass,
is ascribed to the smallness of three different couplings and not just to the tuning of a single
parameter of the Lagrangian. These new couplings are $\lambda, f_q, f$ of the model of eq. (5) and
they all represent new Yukawa interaction in the superpotential among MSSM states and
states beyond the MSSM.

It is extremely interesting to notice that to provide the type-II seesaw model with such a
natural origin for the $\Delta HH$ term it has been necessary to introduce an exotic hypercharge
field. This field has the property that it allows us to write interactions for the triplet $\Delta$ that
extend those of the minimal SUSY type-II seesaw model. These interactions are used to
generate radiatively the mass of the neutrinos, and therefore they are the very core of our
model.

We have analyzed the observable consequences of the new interaction beyond the gener-
ation of neutrino mass. We found two very interesting consequences of this new dynamics.
We have observed that the dynamics of our new quarks is markedly different from that of
minimal models of heavy quarks. In particular we have found significant deviations from
the picture of heavy quarks that decay mostly in Goldstone bosons and SM quarks, \textit{i.e.}
$Q_H \rightarrow q_{SM}Z, h, W$. The several decay patterns and the corresponding estimates of the de-
cay widths are summarized in Table 2 and Table 3. Among all the decays, those mediated by
the coupling $\lambda$ stand out as the most different from the usual $Q_H \rightarrow q_{SM}Z, h, W$ picture. In
fact there are scenarios where the decays of the heavy quark can have very modest leptonic
activity, as for the decays eqs. (19),(20),(21). The lack of leptons in these signature may
significantly impact on the bound on the mass of the existence of heavy quarks, that are
currently searched only in the decay modes $Q_H \rightarrow q_{SM}Z, h, W$. We have seen that in our
model there are also scenarios where the final states are lepton-rich, as for instance in the
case of eq. (22) that can give rise to resonating pairs of leptons of same charge. This signal
is so spectacular that, reinterpreting the searches for doubly charged scalars [28], one could
extend the limits on heavy quarks up to about 1 TeV. Therefore our study gives further
motivation to enlarge the set of final states considered for LHC searches for heavy quarks.
Furthermore our study calls for of a broader interpretation of current searches for SUSY
and other BSM scenarios, which in fact are also probing the dynamics of non-minimal heavy
quarks.

Additionally we have observed that the matter content and interactions of the model
generates a dimension 9 effective operator that mediates the neutrinoless double beta decay.
This fact \textit{per se} is quite expected in models of Majorana neutrinos, as the mass of the
neutrino breaks lepton number and therefore can be used as a source for $0\nu\beta\beta$. However,
we have found that in our model the rate for $0\nu\beta\beta$ is in general only loosely connected to
the generation of the neutrino mass. The degree of disentanglement between the scale of
neutrino mass and the rate of double beta decay is represented by the ratio $f_q \lambda_1/(f_q \lambda_3)$ in
eq(38). This shows that we can have scenarios in which the double beta decay is mostly due
to the short-range exchange of heavy mediators and only sub-dominantly due to long-range
effects from neutrino mass. This hierarchy of the contributions to $0\nu\beta\beta$ can be attained
especially when the couplings of the new heavy quarks exhibit a sort of inverted hierarchy
$f_q \lambda_1/(f_q \lambda_3) \gg 1$, i.e. the third generation SM quarks are less strongly coupled to the new
quarks than what the first generation does. Very excitingly the predicted rate for $0\nu\beta\beta$
can be in the reach of the next round of experiments. In case a signal will be observed, we
have stressed that in principle the LHC and DIS experiments can cross-check the possible observation of $0\nu\beta\beta$. At the LHC the most important diagnostic is the flavor composition of the final states of the heavy quarks decay eqs. (19), (20), (21). At DIS experiments one expects to see a signal in final states with a $W$ boson in association with a hard lepton that has electric charge opposite to that of the beam.

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