Neutrino and Z' phenomenology in an anomaly-free $U(1)$ extension: role of higher-dimensional operators

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ABSTRACT: We consider an anomaly-free $U(1)$ extension of the Standard Model with three right-handed neutrinos (RHNs) and two complex scalars, wherein the charge assignments preclude all tree-level mass terms for the neutrinos. Considering this setup, in turn, to be only a low-energy effective theory, we introduce higher-dimensional terms *a la* Froggatt-Nielsen to naturally generate tiny neutrino masses. One of the RHNs turns out to be very light, thereby constituting the main decay mode for the Z' and hence relaxing the LHC dilepton resonance search constraints. The lightest RHN has a lifetime comparable to or bigger than the age of the Universe, and, hence, could account for a non-negligible fraction of the dark matter.

KEYWORDS: Beyond Standard Model, Neutrino Physics, Gauge Symmetry

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1 Introduction

Of all the lacunae besetting the Standard Model (SM), the existence of Dark Matter (DM) and the flavour problem are rather vexing ones. While the masses of the charged fermions as well as the mixing amongst the quarks can be explained by postulating a certain set of Yukawa couplings, the large hierarchy between these is a rather disquieting feature. Over the last few decades, several disparate sets of theories have been proposed including (but not limited to) (i) additional symmetries, discrete or continuous, gauged or global [1–11], (ii) quark compositeness, whether in terms of further constituents, bound together by some unspecified force [12–15] or, in the more modern parlance, in terms of a higher-dimensional theory, often with a nontrivial gravitational background [16–20]. Very often, though, such

efforts are faced with undesirable phenomenological consequences (unsuppressed flavour-changing neutral currents being one such), and these issues can be cured only through the introduction of further complications. Even more damagingly, no corroborating evidence has been found in terms of additional particles (that many such scenarios posit) or interactions.

The situation has been exacerbated in recent years by the observation of neutrino oscillations [21–25] and these seemingly call out for nonzero neutrino masses. Indeed, the consequent mixing angles in the neutrino sector are quite well-determined and so is one difference in the squares of masses [26, 27]. For the other independent difference, only the magnitude (and not the sign) is known and also unknown are the nontrivial phases that are possible in the mixing matrix. And while the oscillation data is only sensitive to the difference in mass-squareds and not the absolute mass scale, the latter is very-well constrained to $\sum m(\nu_i) < (0.340 - 0.715) \text{ eV}$ [26] — where ν_i are (cosmologically) stable light neutrinos — from a host of cosmological data. Interestingly, direct bounds from terrestrial experiments (such as decays) are significantly weaker [28].

Of course, neutrino masses (absent within what is known as the SM) can be trivially obtained by introducing right-handed neutrino (RHN) fields ν_{jR} and ascribing additional Yukawa terms. Tuning the said couplings so as to obtain the requisite masses and mixings is a seemingly trivial task, once a further hierarchy in the Yukawa couplings (necessitated by the smallness of the neutrino masses) is accepted. The pitch is further queered, though, by the fact that with the RHNs being gauge singlets, terms such $(\overline{\nu_{jR}})^c \nu_{kR}$ are gauge invariant, and being unprotected by any symmetry,¹ can be arbitrarily large. Beset with such large Majorana masses, the ν_{jR} can be integrated out from the low-energy theory, leaving the SM neutrinos with tiny masses. Indeed, this very observation led to a cure in the form of the *seesaw mechanism*, wherein a large (dynamical) scale is set for the Majorana mass matrix m_M , with and the usual Dirac mass matrix (m_D) for the neutrinos being unsuppressed, so that on integrating out the heavy fields, the light fields are left with an effective mass matrix $m_{\text{eff}} \sim m_D^T m_M^{-1} m_D$ which, on diagonalization, should yield the observed masses and mixings.

With m_D being proportional to the vacuum expectation value of the SM Higgs H , the aforementioned structure could have been divined in an effective theory. Written in a gauge-invariant form, the Weinberg operator [29] reads $c_{ij} \overline{L_i^c} L_j H H / \Lambda$ where L_i are the left-handed lepton doublets and Λ is the cutoff scale (perhaps allied to the largest eigenvalue of m_M above). The dimensionless constants c_{ij} constitute a symmetric matrix and can be thought of as parametrizing the structure of $m_D^T m_M^{-1} m_D$. Once again, choosing c_{ij} / Λ to be small enough and ascribing the necessary structure, the correct set of masses and mixings can be obtained.

All of the aforementioned mechanisms to generate the light neutrino masses “suffer” from the requirement of either postulating very tiny couplings or a very large scale (e.g., $m_M \sim 10^{11} \text{ GeV}$ for the so-called “type-I” seesaw mechanism) rendering such theories es-

¹Lepton number conservation is only accidental within the SM and is, actually, broken by nonperturbative effects.

entially untestable by current experiments. It would, thus, be very attractive to have a theory for neutrino masses with a characteristic scale $\mathcal{O}(10\text{ TeV})$ or lower so that it is testable at the LHC, the B -factories etc. Examples are scenarios [30, 31] of TeV scale RHNs with a significant mixing with the SM ν s achieved through the realization of low scale see-saw through non-trivial flavor structure. Similarly, models with radiative neutrino mass generation and/or inverse seesaw [32, 33] also exist; with relatively light RHNs, these can be probed at colliders. For more insight into radiative neutrino mass models one can see the review [34] and the references there in. Looking for well-motivated scenarios that incorporate experimentally testable RHNs in a more broader scheme, is the goal we set for this work.

Before we start this in right earnest, it is worthwhile to remind ourself of a particularly elegant proposal addressing the fermion mass hierarchy. As Froggatt and Nielsen (FN) [35] pointed out, ascribing the Higgs and quark fields with some extra charges (corresponding to a discrete or a continuous symmetry) would, in general, render the usual Yukawa terms untenable. Instead, higher dimension terms could be written by inserting an appropriate number of a “flavon” scalar field $\mathcal{F}$. The choice of the charges would dictate the powers of the ratio $\langle\mathcal{F}\rangle/\Lambda$ (Λ being the cutoff scale) in the effective mass terms and hence their scales. Apparently, then, this simple ruse can lead to correct masses and mixings without the need for imposing a large hierarchy in the Yukawa couplings [36]; and the non-renormalizable nature of the theory could be explained as being the result of integrating out unspecified fields, the nature of which would depend on the particular ultraviolet completion of the FN-scenario. Unfortunately, though, the simplest such models turn out to be phenomenologically inconsistent, failing to satisfy the constraints from rare decays while still explaining the masses and the mixings.

While we would not discuss charged fermion masses in this paper, it is still instructive to examine the FN mechanism and, in particular, where the flavor charges correspond to a U(1) symmetry. An extra gauged U(1) can, of course, appear in many a scenario, ranging from flavor models to theories of compositeness to GUTs [37–57]. Naturally, the phenomenological consequences are very well studied [58–70] and strong constraints (in the mass-gauge coupling plane) emerge from a variety of measurements, ranging from rare decays, anomalous magnetic moments of the electron or muon, electroweak precision tests (performed at the Z -peak) to direct observation at the LHC. The relative strengths of the constraints are determined by the U(1) charge assignments for the light SM fermions. The latter, of course, are not entirely arbitrary as the U(1) extension needs to be anomaly-free. In particular, some of the strongest constraints emanate from the lack of unexplained, yet discernible, peaks in the invariant mass spectra for dijet or dilepton production at the LHC.

In this paper, we pursue a modest goal. Starting with an anomaly-free U(1) extension of the SM (augmented by the mandatory three RHN fields), we employ a Froggatt-Nielsen-like mechanism, but restricted strictly to the neutrino sector. We find that (a) the neutrino masses and mixings can be explained with very moderate choices for the Yukawa couplings and a U(1) scale of a few TeVs; (b) simultaneously, the dilepton branching fraction of the Z' is suppressed so that even with very natural choices of parameters, a Z' as light as 3 TeV is perfectly consistent with the LHC results and, yet, (c) novel signatures are predicted at the LHC.

An additional bonus is the natural emergence of a viable Dark Matter candidate, thereby addressing the second (and, perhaps, even more pressing) lacuna of the SM that we had alluded to. In fact, as many as two ultralight particles (a pseudoscalar and a RHN) appear in the spectrum, with their masses uplifted only by higher-dimensional operators. While not strictly stable, the lighter of these has a lifetime comparable to or greater than the age of the universe. Being charged under the extra $U(1)$ (although neutral under the SM gauge group), their interactions are large enough to be interesting in the context of both cosmology and direct detection.

The rest of the paper is organized as follows. We discuss the anomaly-free $U(1)_z$ extension model in section 2 including gauge, scalar and fermionic sectors of the model. We devote the phenomenology of the neutrino sector and Z' boson in sections 3 and 4 respectively wherein we discuss various exclusion limits. In section 5, we explore the possibility whether the new ultralight particles present in our model can act as a suitable dark matter candidate. Finally, we summarize and conclude in section 6.

2 The $U(1)_z$ extension

The gauge sector of the SM (i.e. $SU(3)_c \times SU(2)_L \times U(1)_Y$) is extended by a new $U(1)_z$ gauge symmetry with the associated gauge coupling g_z . Presence of multiple $U(1)$ s in a gauge theory can, in general, lead to kinetic mixing.² Since, in this paper, we are primarily interested in the effective TeV-scale phenomenology, we make the simplifying assumption that no such kinetic mixing exists.³ We further assume that the SM fields are charged under the new $U(1)_z$ and the corresponding quantum number of a field $\mathcal{F}$ being denoted by $z_{\mathcal{F}}$. With the SM fields too allowed to have nonzero $z_{\mathcal{F}}$, anomaly cancellation is a concern. Postponing this concern until later, we begin by considering the bosonic sector.

2.1 Scalars and symmetry breaking

At a scale much higher than the electroweak symmetry breaking (EWSB) scale, the $U(1)_z$ is broken by one or more SM singlets χ_A carrying charges z_{χ_A} . While a single χ suffices for the requisite symmetry breaking, we generalize the situation for phenomenological reasons which will be clear later. The Lagrangian for the scalar sector is given by

$$\mathcal{L} \supset (D^\mu H)^\dagger (D_\mu H) + \sum_A \left(\tilde{D}^\mu \chi_A \right)^\dagger \left(\tilde{D}_\mu \chi_A \right) - V(\{H^\dagger H\}, \{\chi_A^\dagger \chi_A\}), \quad (2.1)$$

where H denotes the SM Higgs doublet and our assumption that the charges z_{χ_A} are such that trilinear terms in the potential are not admissible would be vindicated later. The covariant derivatives, for a generic field $\mathcal{F}$ is, of course, given by

$$D_\mu = \partial_\mu - ig_s T^a G_\mu^a - ig_w \frac{\sigma^i}{2} W_\mu^i - ig_y \frac{Y}{2} B_\mu - ig_z \frac{z_{\mathcal{F}}}{2} X_\mu; \quad \tilde{D}^\mu = \partial^\mu - ig_z \frac{z_{\mathcal{F}}}{2} X^\mu, \quad (2.2)$$

²While it is always possible to rotate away such a kinetic mixing at one particular scale, such a rotation would be manifested in the gauge-boson mass matrix as well as the couplings with matter. And, even, in the event of a symmetry being unbroken, such a mixing would be regenerated at other scales.

³We have checked explicitly that the mixing engendered by quantum corrections between the scale Λ and the TeV-scale is too small to be of any consequence.

with X_μ being the new gauge boson. The second, third and the fourth terms on the r.h.s. of the first of the above equations correspond to the $SU(3)_c$, $SU(2)_L$ and $U(1)_Y$ gauge groups of the SM with gauge couplings g_s , g_w and g_y respectively.

The symmetry is broken, in two steps, by the vacuum expectation values (vevs) of χ_A and H fields.

$$\langle \chi_A \rangle \equiv \frac{x_A}{\sqrt{2}}, \quad \langle H \rangle \equiv \begin{pmatrix} 0 & \frac{v_h}{\sqrt{2}} \end{pmatrix}^T. \quad (2.3)$$

We assume that, in case multiple χ_A are invoked, the corresponding vevs x_A are of the same order, i.e., no hierarchy is introduced between them. Furthermore, for the sake of simplicity, we do not admit spontaneous CP violation, or in other words, nontrivial phases between the vevs.

While the expression for W -boson mass M_W remains unchanged, the (mass)² matrix for the neutral gauge bosons is now modified to

$$\mathcal{M}^2 = \frac{1}{4} \begin{bmatrix} g_y^2 v_h^2 & -g_y g_w v_h^2 & g_y g_z z_H v_h^2 \\ -g_y g_z v_h^2 & g_w^2 v_h^2 & -g_w g_z z_H v_h^2 \\ g_y g_z z_h v_h^2 & -g_w g_z z_H v_h^2 & g_z^2 \left(z_H^2 v_h^2 + \sum_A z_{\chi_A}^2 x_A^2 \right) \end{bmatrix}. \quad (2.4)$$

Clearly, $\mathcal{M}^2$ is a rank-2 matrix, and can be diagonalized by an orthogonal matrix $\mathcal{O}$ defined through

$$\begin{bmatrix} A_\mu \\ Z_\mu \\ Z'_\mu \end{bmatrix} = \begin{bmatrix} \cos w & \sin w & 0 \\ -\cos t \sin w & \cos t \cos w & \sin t \\ \sin t \sin w & -\sin t \cos w & \cos t \end{bmatrix} \begin{bmatrix} B_\mu \\ W_{3\mu} \\ X_\mu \end{bmatrix} \equiv \mathcal{O}^\dagger \begin{bmatrix} B_\mu \\ W_{3\mu} \\ B'_\mu \end{bmatrix}. \quad (2.5)$$

The Weinberg angle w remains unaltered, namely $w = \tan^{-1}(g_y/g_w)$ whereas for $e = g_w \sin w$, the $Z \leftrightarrow Z'$ mixing angle t is given by

$$\frac{4e z_H g_z}{\sin 2w} \cot 2t = \frac{g_z^2}{v_h^2} \left(\sum_A z_{\chi_A}^2 x_A^2 + z_H^2 v_h^2 \right) - \frac{4e^2}{\sin^2 2w}. \quad (2.6)$$

The heavy neutral gauge boson masses are given by

$$M_{Z,Z'}^2 = \frac{e^2 v_h^2 \cos^2 t}{\sin^2 2w} + \frac{g_z^2}{4} \left(z_H^2 v_h^2 + \sum_A z_{\chi_A}^2 x_A^2 \right) \sin^2 t \mp \frac{e g_z z_H v_h^2}{2 \sin 2w} \sin 2t. \quad (2.7)$$

The shift in M_Z imposes a constraint on the parameter space of the model, which, as far as this sector is concerned, could be considered of being three-dimensional, namely defined by g_z , z_H and the combination $\sum_A z_{\chi_A}^2 x_A^2$. Note, however, that the above tree-level expression cannot be immediately compared with the experimentally measured M_Z as quantum corrections need to be included. We return to this point later.

Another example of such changes would be that wrought by the scalar sector. The very inclusion of a single χ_A field results, after the breaking of the $U(1)_z$ and in the unitary gauge, in an additional scalar field. Although a SM singlet, this can mix with the SM Higgs

field (owing to terms such as $H^\dagger H \chi^\dagger \chi$) resulting in two physical scalars $h_{1,2}$. It thus needs to be ensured that at least one of the two eigenstates has a mass of 125 GeV and couplings (both gauge and Yukawa) not significantly different from the SM Higgs. As it would turn out, this holds almost trivially for the parameter space that we are interested in. It should also be apparent that once one of these (say h_1) is forced to be very SM-like, the other (h_2), being singlet-dominated, would have very small production cross sections at the LHC. The modes of interest would be $q\bar{q} \rightarrow Z' h_2$ (analogous to the Bjorken process) and $gg \rightarrow h_1 h_2$ with (h_2 -sstrahlung). Understandably, such an h_2 would have escaped detection thus far.

The introduction of each additional χ_A results in the physical spectrum being enhanced by a pair of spin-0 particles, one scalar and one pseudoscalar. Particles in each sector could mix amongst themselves with inter-sector mixing disallowed as long as additional CP -violation (explicit or spontaneous) is not introduced in the Higgs sector. It should be noted that this introduces a new class of modes, namely the Z' (on-shell or off-shell) going to a scalar-pseudoscalar pair.

2.2 Example with two χ fields

As an example, consider the special case of there being two such fields χ_1 and χ_2 . This would prove to be of particular interest in the context of neutrino mass generation. If the corresponding $U(1)_z$ charges z_{χ_1} and z_{χ_2} are not integral multiples of each other,⁴ the most general form of the scalar potential invariant under $U(1)_z$ is given by

$$V(\chi_1, \chi_2) = -\mu_1^2 \chi_1^\dagger \chi_1 - \mu_2^2 \chi_2^\dagger \chi_2 + \frac{\lambda_1}{2} (\chi_1^\dagger \chi_1)^2 + \frac{\lambda_2}{2} (\chi_2^\dagger \chi_2)^2 + \lambda_{12} (\chi_1^\dagger \chi_1) (\chi_2^\dagger \chi_2). \quad (2.8)$$

Here, we have deliberately suppressed terms containing the SM Higgs field. This simplifying approximation, apart from being very good at energies far above the electroweak scale, serves to highlight certain salient features. As is immediately apparent, $V(\chi_1, \chi_2)$ is invariant under a global $U(1) \times U(1)$ with each factor associated with one of $\chi_{1,2}$. The most general symmetry breaking would, thus, result in two putative Goldstone fields $\mathcal{G}_{1,2}$ (massless at this scale).

One linear combination of the two $U(1)$ s is gauged, viz. $U(1)_z \sim z_{\chi_1} [U(1)]_1 + z_{\chi_2} [U(1)]_2$, and the corresponding combination of $\mathcal{G}_{1,2}$ appears as the longitudinal component of B'_μ . The orthogonal combination appears in the spectrum as a *massless* pseudoscalar A . While it may seem that the presence of a Goldstone in the theory would render it phenomenologically nonviable, this is not necessarily so, as we argue below.

A true Goldstone would translate to a long-range force. And, even if it acquired a small mass through an explicit symmetry-breaking term, such a particle would still make its presence felt in a variety of interactions. Some of the strongest bounds emanate from low-energy processes including, but not limited to, astrophysical ones. Indeed, Majoron models (i.e., ones, analogously to us, wherein a Goldstone arises in spontaneously breaking a lepton number symmetry in the quest to achieve Majorana masses [71] have been well-studied in this context. As it turns out, though, in the present case, the (could-be) Goldstone has very suppressed couplings with the SM fermions (thanks to the quantum number assignments)

⁴All we really need to ensure is that ratio of the charges be different from 1, 2 or 3.

and, consequently, such bounds, whether from terrestrial experiments [72] or astrophysical considerations [73], are satisfied easily.

While it might seem that, post-EWSB, the SM Higgs might decay into a pair of these Goldstones, note that the shift-symmetry in the potential (or, equivalently, expression for the divergence of the scale current [74, 75]) precludes this unless the Goldstone itself acquires a mass. It should be realised, though, that the Goldstone could, in general, be lifted by a variety of sources including, but not limited to quantum corrections, rendering it a pseudo-Nambu-Goldstone Boson (pNGB). For example, consider an effective theory exemplified by the inclusion of higher-dimensional terms in the Lagrangian parametrizing unknown effects emanating from physics at still higher energies. While eq. (2.8) represents the most general gauge-invariant potential consistent with renormalizability, once nonrenormalizable terms are allowed, more terms can be present. In the next section, we would argue for $z_{\chi_1} = -3/4$, $z_{\chi_2} = -4$, and for such a case the lowest-dimensional term that breaks the global $[U(1)]^2$ down to $U(1)_z$ is $\chi_1^{16}\chi_2^{*3}$. Such a large engineering dimension of the operator would, typically, generate only a very small mass for the pNGB *viz.* $\mathcal{O}(x_i^{16}/\Lambda^{15})$ where Λ is the cutoff scale. For $x_i/\Lambda \lesssim 0.1$, a very reasonable restriction, this would leave the pNGB in a milli-eV range, reminiscent of axionic dark matter models. As we shall see later, this effect would, most likely, be dominated overwhelmingly by other possible upliftings of the pNGB field A . It suffices to say, at this stage, that this immediately vacates the aforementioned no-go theorem and leads, post-EWSB, to the decay of the SM Higgs to a pair of this (SM-singlet) A -particles [74, 75] through an effective trilinear HAA coupling. Indeed, such decays have long been considered for Majoron-like fields [76, 77], to derive constraints on the respective theories. It is interesting to note that the data on the 125 GeV scalar still allows for a non-significant invisible decay of the particle [78] and this can be used to implement an upper bound on these quartic couplings.

It needs to be ascertained, though, whether $\chi_{1,2}$ are allowed to have gauge-invariant Yukawa terms involving any new fermions in the theory (such as RHN fields) and whether these explicitly break $[U(1)]^2$ down to $U(1)_z$. If such be the case, the pNGB would be lifted courtesy loop corrections.⁵

Were it desirable to substantially raise the Goldstone, it could be trivially done through the introduction of a third singlet scalar χ_3 , which, in principle, could actually increase the global symmetry to $[U(1)]^3$. On the other hand, this may allow for renormalizable terms breaking down the global symmetry to just a single $U(1)$ to be identified with the gauged $U(1)_z$. For example, working with the previously assigned quantum numbers for $\chi_{1,2}$, if one introduces a third scalar χ_3 with a charge $13/4$, then a term such as $(\mu_{123}\chi_1^*\chi_2\chi_3 + \text{H.c.})$ would break the symmetry softly. On the other hand, if $z_{\chi_3} = 5/2$, then a hard breaking is achieved through $(\lambda_{1123}\chi_1^2\chi_2\chi_3^* + \text{H.c.})$.

It is easy to see most such augmentation of the scalar sector does not materially alter low energy phenomenology except, perhaps, to ameliorate some issues with the evolution history of the early universe. As far as LHC signals go, the most drastic changes would

⁵It should be obvious that, with the couplings of $\chi_{1,2}$ with H or B'_μ preserving $[U(1)]^2$, the corresponding corrections to $V(\chi_1, \chi_2)$ would not lift the Goldstone.

be in the decays of the Z' wrought by the proliferation of states (three and two SU(2)-singlet scalars and pseudoscalars). Since these effects are easily computed and are not very germane to the issues that we are primarily interested in, we will not discuss such a three-singlet scenario any further.

To start with let us work under the assumption that the χ -sector is essentially decoupled from the SM Higgs sector. This is not too drastic an approximation at energy scales much higher than the electroweak scale (except as far as the decays of the χ_i into the SM Higgs is concerned). The potential, then, is described by that in eq. (2.8). Denoting the χ fields, post symmetry breaking, by

$$\chi_{1,2} = \frac{1}{\sqrt{2}} (x_{1,2} + \xi_{1,2} + i\rho_{1,2}) , \quad (2.9)$$

where $\xi_{1,2}, \rho_{1,2}$ are real fields, the massless pseudoscalar is given by

$$A = \rho_1 \sin \gamma_A - \rho_2 \cos \gamma_A , \quad \tan \gamma_A = \frac{z_{\chi_2} x_2}{z_{\chi_1} x_1} , \quad (2.10)$$

with the orthogonal combination being absorbed to reappear as the longitudinal mode⁶ of the Z' . The mass-squared matrix for the two scalars $\xi_{1,2}$ reads

$$M_{\xi_i}^2 = \begin{pmatrix} \lambda_1 x_1^2 & \lambda_{12} x_1 x_2 \\ \lambda_{12} x_1 x_2 & \lambda_1 x_1^2 \end{pmatrix} ,$$

leading to mass eigenstates $\tilde{\xi}_{1,2}$ defined by

$$\begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = \begin{pmatrix} \cos \alpha_\chi & \sin \alpha_\chi \\ -\sin \alpha_\chi & \cos \alpha_\chi \end{pmatrix} \begin{pmatrix} \xi_1 \\ \xi_2 \end{pmatrix} , \quad \tan(2\alpha_\chi) = \frac{2\lambda_{12} x_1 x_2}{\lambda_1 x_1^2 - \lambda_2 x_2^2} , \quad (2.11)$$

with the corresponding masses being

$$M_{H_1, H_2}^2 = \frac{1}{2} [\lambda_1 x_1^2 + \lambda_2 x_2^2 \pm |\lambda_1 x_1^2 - \lambda_2 x_2^2| \sec(2\alpha_\chi)] . \quad (2.12)$$

2.3 Fermionic sector and anomalies

Since one of our primary goals is to explain neutrino masses and mixings, we must include extra neutral fermions and at least two of them. This is the only addition we propose in this sector, and as we would shortly see, invoking three such right-handed fields is not only enough to ensure the cancellation of all possible anomalies,⁷ but also leads to very interesting phenomenological consequences.

For the sake of simplicity, we consider the U(1)_z charges to be family-blind, as far as the SM fermions are concerned, denoting these to be z_Q (quark doublets), z_L (lepton

⁶With the SM Higgs not yet acquiring a nonzero vacuum expectation value, there is no Z - Z' mixing at this stage.

⁷While, gauge anomalies can also be canceled using the Green-Schwarz mechanism — we refer the reader to ref. [79] for a phenomenological discussion of such constructions — we eschew this in favour of a more canonical approach.

doublets), z_u, z_d (the right-handed up-like and down-like quarks respectively) and z_e for the right-handed charged leptons. Similarly, three RHN fields N_i are assigned charges z_i , not necessarily equal. Before we consider the fermion masses and, thereby, relate these to z_H and z_{χ_A} , let us first discuss the anomalies defined as $\mathcal{A} \equiv \text{tr}_L(T_a T_b T_c) - \text{tr}_R(T_a T_b T_c)$ where T_a are the symmetry generators and the traces are over left- and right-handed fermions. The SM gauge anomalies, of course, remain unaltered and the only nontrivial quantities are those pertaining to $U(1)_z$ and are listed below

Anomaly	Expression
$[\text{SU}(3)_c]^2 U(1)_z$	$2z_Q = z_u + z_d$
$[\text{SU}(2)]^2 U(1)_z$	$3z_Q + z_L = 0$
$[\text{U}(1)_Y]^2 U(1)_z$	$z_Q + 3z_L = 8z_u + 2z_d + 6z_e$
$U(1)_Y [U(1)_z]^2$	$z_Q^2 - z_L^2 = 2z_u^2 - z_d^2 - z_e^2$
$[U(1)_z]^3$	$6z_Q^3 + 2z_L^3 = 3z_u^3 + 3z_d^3 + z_e^3 + \sum_{i=1}^3 z_i^3$

It is easy to see that, using the first two conditions, the third simplifies to $2z_Q + z_u + z_e = 0$. Similarly, the first three, together, imply that the fourth one is satisfied identically. And, finally, the fifth one simplifies to $\sum_{i=1}^3 z_i^3 = 3(z_u - 4z_Q)^3$. It is also easy to see that the mixed gauge-gravity anomaly ($R^2 U(1)_z$) does not present an independent constraint.

2.4 $U(1)_z$ charge assignment

The existence of mass terms for the charged fermions demands that

$$z_H = z_L - z_e = z_Q - z_d = z_u - z_Q. \quad (2.13)$$

Note that only one of these equations is independent once anomaly cancellations have been imposed (in fact, just the $[\text{SU}(3)_c]^2 U(1)_z$ and $[\text{SU}(2)]^2 U(1)_z$ are enough). The $U(1)_z$ charges of the SM fields can, then, be expressed in terms of just two parameters, say z_u and z_Q . Note, however, that, for any $U(1)$ theory, one combination of charges can always be taken to be unity, without any loss of generality. In the present case, we shall choose $z_u - 4z_Q = 1$ and consider z_Q to be the remaining free parameter. The consequent charge assignments have been displayed in table. 1.

The charges z_i for the N_{iR} fields, thus, need to satisfy

$$\sum_i z_i^3 = 3,$$

and the solution space is a two-dimensional one. Restricting ourselves to rational values, the simplest assignment would be $z_i = 1$, a choice that has been explored in a different context [80]. This, though, is unsuitable as far as neutrino mass generation is concerned. Consequently, we adopt the next simplest choice, namely $z_{1,2} = 4, z_3 = -5$.

It might be argued that, unlike in the example of ref. [80], such a non-universal choice is not amenable to the simplest grand unified theories. However, identifying $U(1)_z$ to be

	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
q_L	3	2	1/6	z_Q
u_R	3	1	2/3	$1 + 4z_Q$
d_R	3	1	$-1/3$	$-1 - 2z_Q$
ℓ_L	1	2	$-1/2$	$-3z_Q$
e_R	1	1	-1	$-1 - 6z_Q$
H	1	2	1/2	$1 + 3z_Q$
N_{1R}, N_{2R}	1	1	0	4
N_{3R}	1	1	0	-5
χ_1	1	1	0	z_{χ_1}
χ_2	1	1	0	z_{χ_2}

Table 1. The charge assignments for the fermions and scalars of the model.

a part of a GUT entails strong constraints, not the least being the those imposed by the requirement of gauge coupling unification on the one hand and the suppression of proton decay on the other. On the other hand, nonminimal models (such as those arising from orbifoldings) incorporating extra (incomplete) multiplets of, say E_6 , can, in principle, be built incorporating such fields. While such extended flavon models can address many issues, including inflation [81], it is not our intention to attempt this. Our goal is more modest, namely to provide for an understanding of neutrino masses as well as to offer a viable candidate for the Dark Matter.

3 Neutrino masses

While the Yukawa (and, hence, the mass) terms for the charged fermions proceed as within the SM, *viz.*,

$$\mathcal{L}_{\text{Yuk.}} = y_{ij}^u \bar{Q}_{Li} u_{Rj} \tilde{H} + y_{ij}^d \bar{Q}_{Li} d_{Rj} H + y_{ij}^e \bar{L}_{Li} e_{Rj} H + \text{H.c.},$$

where $\tilde{H} = i\sigma_2 H^*$, dimension-4 gauge invariant Yukawa (or even bare mass) terms are not possible for the neutrinos. The situation changes if the theory is treated not as a fundamental one, but only as the low-energy limit of some more fundamental theory operative at some scale Λ or higher. Freed of the restriction of being renormalizable, the effective field theory would admit higher-dimensional terms of the form⁸

$$\begin{aligned} \mathcal{L}_{\nu\text{mass}} &= \mathcal{L}_{\text{Dirac}} + \mathcal{L}_{\text{Wein.}}; \\ \mathcal{L}_{\text{Dirac}} &= \sum_{i=1}^3 \sum_{\alpha=1}^2 y_{i\alpha} \bar{L}_{iL} N_{\alpha R} \tilde{H} \frac{\chi_1^{a_1} \chi_2^{a_2}}{\Lambda^{|a_1|+|a_2|}} + \sum_{i=1}^3 \tilde{y}_i \bar{L}_{iL} N_{3R} \tilde{H} \frac{\chi_1^{a_3} \chi_2^{a_4}}{\Lambda^{|a_3|+|a_4|}} + \text{H.c.}; \end{aligned}$$

⁸While analogous terms can be written for the charged fermions as well, these would be subdominant to the usual Yukawa terms and we omit all discussions thereof.

$$\begin{aligned} \mathcal{L}_{\text{Wein.}} = & \sum_{i,j=1}^3 w_{ij} \overline{L_{iL}} L_{jL} H H \frac{\chi_1^{b_1} \chi_2^{b_2}}{\Lambda^{|b_1|+|b_2|+1}} + \sum_{\alpha,\beta=1}^2 s_{\alpha\beta} \overline{N_{\alpha R}} N_{\beta R} \frac{\chi_1^{b_3} \chi_2^{b_4}}{\Lambda^{|b_3|+|b_4|-1}} \\ & + \sum_{\alpha=1}^2 s_{\alpha 3} \overline{N_{\alpha R}} N_{3R} \frac{\chi_1^{b_5} \chi_2^{b_6}}{\Lambda^{|b_5|+|b_6|-1}} + s_{33} \overline{N_{3R}} N_{3R} \frac{\chi_1^{b_7} \chi_2^{b_8}}{\Lambda^{|b_7|+|b_8|-1}} + \text{H.c.}, \end{aligned} \quad (3.1)$$

where the couplings $y_{i\alpha}, \tilde{y}_i, w_{ij}, s_{\alpha\beta}, s_{\alpha 3}$ and s_{33} are dimensionless and the exponents satisfy

$$\begin{aligned} z_{\chi_1} a_1 + z_{\chi_2} a_2 &= -3 & z_{\chi_1} a_3 + z_{\chi_2} a_4 &= 6 \\ z_{\chi_1} b_1 + z_{\chi_2} b_2 &= -2 & z_{\chi_1} b_3 + z_{\chi_2} b_4 &= -8 \\ z_{\chi_1} b_5 + z_{\chi_2} b_6 &= 1 & z_{\chi_1} b_7 + z_{\chi_2} b_8 &= 10. \end{aligned} \quad (3.2)$$

It should be realized that only integer solutions for the exponents are permissible as non-integral values would imply nonlocal operators. Negative values for the exponents are to be interpreted as positive powers of $\chi_1^* (\chi_2^*)$ as the case may be.

Before discussing the ramifications of $\mathcal{L}_{\nu\text{mass}}$, it is instructive to remind ourselves of the possible origin of the same. As can be readily recognized, these can arise from a UV-complete theory once a slew of fields (especially fermionic ones) are integrated out. Clearly, these fields must have masses larger than Λ and carry nonzero U(1) charges. It might be argued, then, that the requirement of the effective theory being anomaly-free is a superfluous one, for the anomaly(ies), being a child of the UV regularization, could, in principle, be canceled by the heavy fermions Ψ_i . However, note that the Ψ_i themselves should be vector-like, as else their masses can only arise from spontaneous breaking of the U(1) symmetry and, thus, should be below Λ . And since vector-like fermions do not contribute to gauge anomalies, the effective theory better be anomaly-free.

A further issue pertains to the relative strengths of tree-order and loop-level contributions in the generalization of the FN-mechanism that our theory really represents. Consider a typical term in $\mathcal{L}_{\nu\text{mass}}$ which has n powers of, say, χ_1 . Letting the χ_1 lines go into the vacuum (courtesy spontaneous symmetry breaking) gives us a factor of $(x_1/\Lambda)^n$. On the other hand, closing a χ_1 loop would, typically, give us a factor of $\mathcal{O}(\Lambda^2/16\pi^2 x_1^2)$, as the loop momentum would need to be cut off at the scale Λ . Thus, the exclusion of loops is well motivated for $\Lambda > x_j \gtrsim \Lambda/(4\pi)$, inequalities that we would satisfy in further calculations.

With each of the terms in $\mathcal{L}_{\text{Wein.}}$ violating lepton-number, it is tempting to characterize the corresponding mass terms, realised on breaking the $\text{U}(1)_z$ symmetry, as Majorana masses. An argument against this would be the fact that, in the canonical sense, a Majorana particle may not have any nonzero additive quantum numbers, whereas each of ν_{iL} and N_{iR} certainly do. Rather, these terms should be thought of as the generalization of the Weinberg-operator [29] that is allowed as a dimension-5 correction to the SM. Indeed, on the breaking of the $\text{U}(1)_z$ symmetry, the corresponding charge is no longer a valid quantum number in the ensuing theory, and the mass term generated thereupon can indeed be thought of as a Majorana mass.

As a very specific case, let us consider the assignment

$$z_{\chi_1} = -3/4, \quad z_{\chi_2} = -4, \quad (3.3)$$

which leads to rather interesting phenomenology.⁹ With this choice, the masses for $N_{1,2}$ sub-sector are relatively unsuppressed, while terms connecting N_{3R} are highly suppressed. Indeed, retaining just the least suppressed terms in each sector would lead to

$$\begin{aligned}\mathcal{L}_{\text{Dirac}} &= \sum_{i=1}^3 \sum_{\alpha=1}^2 y_{i\alpha} \bar{L}_{iL} N_{\alpha R} \tilde{H} \frac{\chi_1^4}{\Lambda^4} + \sum_{i=1}^3 \tilde{y}_i \bar{L}_{iL} N_{3R} \tilde{H} \frac{\chi_1^{*8}}{\Lambda^8} + \text{H.c.}, \\ \mathcal{L}_{\text{Wein.}} &= \sum_{i,j=1}^3 w_{ij} \bar{L}_{iL}^c L_{jL} H H \frac{\chi_1^8 \chi_2^*}{\Lambda^{10}} + \sum_{\alpha,\beta=1}^2 s_{\alpha\beta} \bar{N}_{\alpha R}^c N_{\beta R} \frac{\chi_2^2}{\Lambda} \\ &\quad + \sum_{\alpha=1}^2 s_{\alpha 3} \bar{N}_{\alpha R}^c N_{3R} \frac{\chi_1^4 \chi_2^*}{\Lambda^4} + s_{33} \bar{N}_{3R}^c N_{3R} \frac{\chi_1^8 \chi_2^{*4}}{\Lambda^{11}} + \text{H.c.}.\end{aligned}\tag{3.4}$$

The terms corresponding to w_{ij} and s_{33} are too small to be of any consequence, and, formally, could be dropped altogether if we restrict ourselves to operators of mass dimension 12 or less. This leads to

$$\begin{aligned}\mathcal{L}_{\nu\text{mass}} &\approx \mathcal{L}^{(5)} + \mathcal{L}^{(8)} + \mathcal{L}^{(12)} + \text{H.c.} \\ \mathcal{L}^{(5)} &\equiv \sum_{\alpha,\beta=1}^2 s_{\alpha\beta} \bar{N}_{\alpha R}^c N_{\beta R} \frac{x_2^2}{\Lambda}, \\ \mathcal{L}^{(8)} &\equiv \sum_{i=1}^3 \sum_{\alpha=1}^2 y_{i\alpha} \bar{L}_{iL} N_{\alpha R} \tilde{H} \frac{x_1^4}{\Lambda^4} + \sum_{\alpha=1}^2 s_{\alpha 3} \bar{N}_{\alpha R}^c N_{3R} \frac{x_1^4 x_2^*}{\Lambda^4}, \\ \mathcal{L}^{(12)} &\equiv \sum_{i=1}^3 \tilde{y}_i \bar{L}_{iL} N_{3R} \tilde{H} \frac{x_1^{*8}}{\Lambda^8},\end{aligned}\tag{3.5}$$

and this is the form that we would be working with henceforth.

3.1 Identifying the mass eigenstates

The neutrino mass matrix, as given in eq. (3.5) can be represented, in the (ν_j, N_3, N_1, N_2) basis, by

$$\mathcal{M}_\nu = \begin{pmatrix} 0_{3 \times 3} & \mathcal{D} \\ D^T & M_N \end{pmatrix}.\tag{3.6}$$

Denoting $\xi \equiv x/\Lambda$, where x is either of $x_{1,2}$ (we assume that there is no large hierarchy between the x_i), the matrices above have structures

$$\mathcal{D} \approx v \xi^4 \begin{pmatrix} y_{13} \xi^4 & y_{11} & y_{12} \\ y_{23} \xi^4 & y_{21} & y_{22} \\ y_{33} \xi^4 & y_{31} & y_{32} \end{pmatrix} \quad M_N \sim \frac{x^2}{\Lambda} \begin{pmatrix} 0 & s_{31} \xi^3 & s_{32} \xi^3 \\ s_{31} \xi^3 & a_1 & 0 \\ s_{32} \xi^3 & 0 & a_2 \end{pmatrix}.\tag{3.7}$$

Note that the N_1 – N_2 sub-sector of the mass matrix can be chosen to be diagonal without any loss of generality, with a_1 and a_2 as the coefficients.

⁹It should be realised that this choice is not a special one and qualitatively similar results would be obtained for many other choices.

While two eigenvalues of M_N are large, *viz.* $\mathcal{O}(x^2/\Lambda)$, these heavy neutrinos still tend to be lighter than the Z' , owing to the smallness of ξ (in comparison to g_X). The third eigenvalue of M_N is much smaller, namely only $\lambda_3 \sim \mathcal{O}(\xi^6 x^2/\Lambda)$. This state has only a small mixing with the heavier ones, with the mixing angles being $\mathcal{O}(\xi^4)$.

More importantly, it might seem that a straightforward application of the seesaw mechanism may not be possible, given that one of the eigenvalues of M_N is smaller than some of the Dirac masses. However, note that the elements $\mathcal{D}_{i1}$ are smaller than the corresponding fulcrum λ_3 by at least a factor of $v\xi/x$. Consequently, the seesaw mechanism goes through trivially, and even the first-order estimate is rather accurate. This is corroborated by a full numerical calculation of the eigensystem of the full 6×6 mass matrix.

On block-diagonalization, we have the effective mass matrix for the light-sector to be given by

$$M_{3 \times 3} = -\mathcal{D} M_N^{-1} \mathcal{D}^T + \mathcal{O}(M_N^{-2}). \quad (3.8)$$

Explicitly,

$$M_{3 \times 3} = \frac{-v^2 \xi^6}{2\Lambda} \begin{pmatrix} M_{11} & M_{12} & M_{13} \\ M_{12} & M_{22} & M_{23} \\ M_{13} & M_{23} & M_{33} \end{pmatrix}, \quad (3.9)$$

where,

$$\begin{aligned} M_{11} &= (y_{11} - y_{12})^2 + 2\xi(y_{11} + y_{12})y_{13} - \xi^2 y_{13}^2 \\ M_{12} &= (y_{11} - y_{12})(y_{21} - y_{22}) - \xi^2 y_{13} y_{23} + \xi(y_{13}(y_{21} + y_{22}) + (y_{11} + y_{12})y_{23}) \\ M_{13} &= (y_{11} - y_{12})(y_{31} - y_{32}) - \xi^2 y_{13} y_{33} + \xi(y_{13}(y_{31} + y_{32}) + (y_{11} + y_{12})y_{33}) \\ M_{22} &= (y_{21} - y_{22})^2 + 2\xi(y_{21} + y_{22})y_{23} - \xi^2 y_{23}^2 \\ M_{23} &= (y_{21} - y_{22})(y_{31} - y_{32}) - \xi^2 y_{23} y_{33} + \xi(y_{23}(y_{31} + y_{32}) + (y_{21} + y_{22})y_{33}) \\ M_{33} &= (y_{31} - y_{32})^2 + 2\xi(y_{31} + y_{32})y_{33} - \xi^2 y_{33}^2 \end{aligned} \quad (3.10)$$

Several points demand attention:

- As argued above, dropping the higher order terms in $M_{3 \times 3}$ is an excellent approximation. Similarly, the mixing between the ν_i and N_α is quite small and the eigenstates of $M_{3 \times 3}$ are predominantly doublets.
- If the Yukawa couplings are real (i.e., if there is no CP violation in this sector), then the real symmetric matrix $M_{3 \times 3}$ is given in terms of 6 parameters. These can be uniquely determined in terms of the three light neutrino masses and the three mixing angles.
- *A priori*, the matrix $\mathcal{D}$ is defined by nine parameters (assumed to be real), and further parameters appear in M_N . However, only six independent combinations may enter $M_{3 \times 3}$, with the rest serving only to define the heavy sector and the tiny mixing between the heavy and light sectors.
- Finally, while exploring the parameter space, it is important to ensure that $(M_{3 \times 3})_{ee}$ is bound by the non-observation of neutrinoless double beta decay [82].

	$\Delta_{21}(\text{eV}^2)$	$\Delta_{32}(\text{eV}^2)$	$\sin^2 \theta_{23}$	$\sin^2 \theta_{13}$	$\sin^2 \theta_{12}$
NH	$7.39_{-0.20}^{+0.21} \times 10^{-5}$	$2.449_{-0.030}^{+0.032} \times 10^{-3}$	$0.558_{-0.33}^{+0.20}$	$2.241_{-0.65}^{+0.66} \times 10^{-2}$	$0.31_{-0.12}^{+0.13}$
IH	$7.39_{-0.20}^{+0.21} \times 10^{-5}$	$-2.509_{-0.032}^{+0.032} \times 10^{-3}$	$0.563_{-0.26}^{+0.19}$	$2.261_{-0.64}^{+0.67} \times 10^{-2}$	$0.31_{-0.12}^{+0.13}$

Table 2. Bounds on neutrino mass and mixing parameters from neutrino oscillation experiments, for both normal hierarchy (NH) and inverted hierarchy (IH) of the neutrino mass spectrum. The error bars shown correspond to 1σ [27].

3.2 Neutrino results

Since the heavy-light mixing is tiny, low-energy experiments are well-described in terms of the light neutrinos alone. To this end, we diagonalize the light part of the mass matrix through

$$U^T M_{3 \times 3} U = \text{diag}(m_1, m_2, m_3), \quad (3.11)$$

where the PMNS matrix U is given by

$$U^T = R_{12}(\theta_{12}) R_{13}(\theta_{13}) R_{23}(\theta_{23}),$$

with R_{jk} denoting a rotation in the jk plane through an angle θ_{jk} . In other words,

$$U = \begin{pmatrix} c_{12}c_{13} & c_{13}s_{12} & s_{13} \\ -s_{12}c_{23} - c_{12}s_{13}s_{23} & c_{12}c_{23} - s_{12}s_{13}s_{23} & c_{13}s_{23} \\ s_{12}s_{23} - c_{12}c_{23}s_{13} & -c_{12}s_{23} - c_{23}s_{12}s_{13} & c_{13}c_{23} \end{pmatrix}, \quad (3.12)$$

where $c_{ij} \equiv \cos \theta_{ij}$ and $s_{ij} \equiv \sin \theta_{ij}$.

The different Yukawa parameters are varied such that the mass squared differences of neutrino mass eigenstates namely $m_i, i = 1, 2, 3$ defined as,

$$\Delta_{21} = m_2^2 - m_1^2$$

and

$$\Delta_{32} = \pm(m_3^2 - m_2^2)$$

along with three mixing angles $\theta_{12}, \theta_{13}, \theta_{23}$ conform to the experimental results as listed in table. 2. Another important constraint which is to be satisfied is cosmological one on the sum of the masses of the light stable neutrinos, namely, [83]:

$$\sum_i m_i < (0.340 - 0.715) \text{ eV}, \quad (3.13)$$

with 95% confidence level (CL). A more stringent limit can be obtained if we take into account Baryonic Acoustic Oscillations [84, 85]. But we have stuck to the limit above.

Clearly, we have more parameters in the theory than there are data in the neutrino sector. Consequently, some parameters are indeterminable, and must be fixed by hand. To ease the task of identifying the crucial dependence, we choose to make some simplifying assumptions:

- While we had already assumed that there was no large hierarchy between the vacuum expectation values for the two new scalars, we eliminate any choice and consider only $x_1 = x_2 = 10 \text{ TeV}$.

- The cutoff scale Λ is fixed at 100 TeV; in other words, $\xi = 0.10$.
- The two heavy RHNs, predominantly $N_{1,2}$, should have a mass $\mathcal{O}(\xi x)$ and these are held at 1.2 TeV and 1.25 TeV respectively. The splitting between them is not germane to the discussion at hand, and has been incorporated just to ensure that the numerical results are not affected by degeneracy.
- The third RHN (N_3 -like) is held at 6 keV. This can be achieved, for example, if $s_{\alpha 3} \approx 0.05$, in eq. (3.4). As it turns out, certain results are quite dependent on this (although, not the light neutrino phenomenology) and we shall return to this in a later section.

It should be realized that the above choices are not special in any way and do not leave qualitative impact on the determination of the rest of the parameter space. These are to be fixed by an analysis of the neutrino oscillation results. Clearly, this last bit would be dependent on the two different neutrino mass hierarchies that are experimentally viable, and we consider them in turn. In doing this, it needs to be borne in mind that, for a very large part of the parameter space, one of the three SM-like neutrinos is distinctly lighter than the others. This is but a consequence of the fact that, were N_3 to be decoupled, the light mass matrix $M_{3 \times 3}$ (resulting from the seesaw mechanism) would be a rank-2 one.

3.2.1 Normal hierarchy

Terming the lightest mass $m_1 \sim 0$, the two other mass eigenvalues would be $m_2 \sim \sqrt{\Delta_{21}}$ and $m_3 \sim \sqrt{\Delta_{32} + \Delta_{21}}$, thereby satisfying the hierarchy $m_1 < m_2 \ll m_3$.

As is well known, the requirement of $\theta_{23} \sim 45^\circ$ imposes strong constraints on the neutrino mass matrix, and is often sought to be explained by family symmetries. In the present context, this is to be ensured by judiciously choosing the couplings y_{ij} guided by eq. (3.10). While many different solutions are possible, given that ξ is small, if we want to eschew very large hierarchies in the couplings and/or large cancellations, we should look for the possibility that $y_{21} - y_{22} \approx y_{31} - y_{32} \sim 0.01$. The first approximate equality ensures that the leading contributions are of the same order, while the second one ensures that cancellations are not extreme (since the couplings themselves would turn out to be of a similar size). The correct value for the heaviest neutrino mass m_3 is obtained for a somewhat larger value of y_{33} , namely $y_{33} \sim \mathcal{O}(0.1)$. As figure 1 shows, it is possible to satisfy all the neutrino constraints in our framework with these Yukawas, without resorting to very small values as is often the case for general seesaw models. Indeed, relatively larger values of the Yukawas $y_{21} \sim 0.08, y_{22} \sim 0.03$ as shown in figure 1 play a pivotal role to have a significant mixing in the 1 – 2 sector. It should be realized that figure 1 does not reflect the entire viable parameter space. The very structure of eq. (3.9), for example, stipulates that the interchange $(y_{11}, y_{21}) \leftrightarrow (y_{12}, y_{22})$ would result in identical masses and mixings. Rather, the parameter space displayed in figure 1 should be considered a representative set of solutions.

Figure 1. Correlation of Yukawa couplings in the Dirac sector for neutrino masses in normal hierarchy. Allowed points after diagonalization of neutrino mass matrix satisfying the bound on total mass of three neutrino species (in yellow), points with satisfying the bound on Δm_{32}^2 (in purple) and allowed points after another bound of Δm_{12}^2 (in red).

3.2.2 Inverted hierarchy

Similar to the normal hierarchy, we have taken a simplified scenario where m_3 is now taken to be zero. Other two mass eigenvalues in this case then become $m_1 \sim \sqrt{\Delta_{23} - \Delta_{21}}$ and $m_2 \sim \sqrt{\Delta_{23}}$. This satisfies the hierarchy $m_3 \ll m_1 \approx m_2$.

Similar in the previous case, some initial assumption on Yukawa couplings are taken such as $y_{11} = y_{21} + 0.01$ and $y_{12} = y_{22} + 0.01$, such that we have similar contribution from $(y_{21} - y_{22})$ and $(y_{11} - y_{12})$. This is required to satisfy the mass of m_1 to be non-zero along with the significant mixing with $\theta_{12} \sim 33^\circ$, in the 1 – 2 sector. There are two more distinct regions of Yukawas with magnitude of similar order to those shown in the plots where all the neutrino constraints are satisfied. In this case also, the allowed Yukawas values satisfying the neutrino results are significant enough, as presented in the figure 2. In the allowed parameter space, the difference $y_{31} - y_{32}$ is relatively larger than the difference $y_{21} - y_{22}$, as depicted in the figure 2, helps to have a negligibly small m_3 term along with significant mixing in the form of $\theta_{23} \sim 45^\circ$.

4 Z' phenomenology

Until now, we have investigated the parameter space of the model only in the context of the light neutrino masses and mixings. As is evident, these observables are sensitive primarily to the Yukawa couplings (the Wilson coefficients in eq. (3.5)). While there is a dependence on x_i and Λ , these appear in a trivial fashion, and as far as the limited number

Figure 2. Correlation of Yukawa couplings in the Dirac sector for an inverted hierarchy of neutrino masses.

of observables available at low energies are concerned, these dependencies can be entirely subsumed in the Wilson coefficients.¹⁰ And, finally, this sector carries virtually no imprint of either the gauge coupling g_z , nor the parameters in the Higgs potential.

The aforementioned parameters are best investigated at colliders, either through direct production or by effecting precision studies. Seven additional physical states now appear: a massive neutral gauge boson Z' , two massive scalars H_1 and H_2 , a relatively light pseudoscalar A , two massive (predominantly right-handed) neutrinos $N_{1,2}$ and, finally, a light neutrino N_3 (again, predominantly right-handed). As for the parameters in the gauge sector, apart from g_z and x_i , we also have the hitherto unfixed charge z_Q (see table 1), in terms of which the $U(1)_z$ charges of all the SM fields were specified, courtesy the requirements of anomaly cancellations. That z_Q is still free is but a reflection of the fact that $U(1)$ charges are, intrinsically, not quantized unless the symmetry had descended from a bigger group.

With the Z' having a substantial coupling to all the SM fields,¹¹ production of the Z' at, say, the LHC would be expected to constitute a sensitive probe of the scenario. Interesting signatures of heavy neutrinos of our model can also be searched for at future lepton colliders [86]. Before delving into the details of the Z' phenomenology, it is amusing to note that choosing z_Q , hitherto a free parameter, appropriately would minimize the production cross-section without affecting the neutrino phenomenology. For example, to

¹⁰In other words, by trivially rescaling the Wilson coefficients, one can reproduce the observed masses and couplings for different sets of (x_i, Λ) .

¹¹Note that while its coupling with any one set of the SM fermions can be switched off entirely, this cannot be done simultaneously for all of them.

the leading order,

$$\sigma(pp \rightarrow Z' + X) \propto (z_q^2 + z_u^2)F_u + (z_q^2 + z_d^2)F_d. \quad (4.1)$$

Here, X symbolizes the rest of the hadronic byproducts and the flux F_u is given by

$$F_u \equiv \int_{M_{Z'}/s}^1 \frac{dx}{x} \left[f_u(x, Q^2) f_{\bar{u}}\left(\frac{M_{Z'}^2}{sx}, Q^2\right) + f_{\bar{u}}(x, Q^2) f\left(\frac{M_{Z'}^2}{sx}, Q^2\right) \right], \quad (4.2)$$

with $\sqrt{s}$ being the total center-of-mass energy available at the LHC, and $f_j(x, Q^2)$ the density of the j^{th} parton for a given momentum fraction x and computed at the scale Q . An analogous expression holds for F_d . While the ratio of the two fluxes (those for the other quark flavours, being subdominant, result in only a very minor correction) is a function of $M_{Z'}^2$, it is intriguing to note that, for most of the region of interest, $\sigma(pp \rightarrow Z' + X)$ is minimized for $z_Q \sim -1/4$. It is interesting to note that a choice of $z_Q = -1/4$ is commensurate (apropos $U(1)_z$ charge quantization) with our assignments for the z_{χ_i} . Consequently, we use this value of z_Q for much of the numerical calculations below. However, this choice is not a very crucial one and somewhat removed values of z_Q might also be chosen as illustrated by an alternative choice for some of the figures.

4.1 Branching ratios

At a collider, of all the new states, the production of the Z' is the easiest. Consequently, we begin by considering its branching fractions. Owing to its large mass, the decay is dominated by the two-body modes and these can be calculated trivially. We analytically compute various partial widths of two-body decay modes of the Z' and the expressions thereof are given in appendix A. Our numerical results are cross-checked against MADGRAPH [87].

As a particular benchmark point in the parameter space, we consider

$$M_{Z'} = 3\text{TeV}, \quad M_{N_1}, M_{N_2} \sim 1\text{TeV}, \quad M_{H_1}, M_{H_2} \sim 1\text{TeV}. \quad (4.3)$$

These values are chosen so as to ensure that all possible two-body decay modes are kinematically open. In figures 3a and 3b, we display the BRs of various decay modes of Z' as a function of $M_{Z'}$ and for two particular values of z_Q . The kinks in these figures are easily understood as manifestations of new thresholds opening up. Inclusion of off-shell (three-body or even four-body final states) serves to smoothen out the kinks.

Given the large $U(1)_z$ charge of the N_3 and its small mass, it is understandable that the $Z' \rightarrow \bar{N}_3 N_3$ mode overwhelmingly dominates. While, individually, the $Z' \rightarrow \bar{N}_{1,2} N_{1,2}$ are expected to be suppressed (compared to the $\bar{N}_3 N_3$ one) by only a factor of $16/25$, these also suffer an additional kinematic suppression, especially at lower values of $M_{Z'}$. For high Z' masses, together, these would slightly overcome the former mode. Of the decays into exclusively SM decay modes, of particular interest (since the corresponding SM backgrounds are not too large) are the ones to $W^+ W^-$ and $Z H_{\text{SM}}$. Equal on account of the Goldstone equivalence theorem, these tend to be small as their amplitudes are proportional to the $Z \leftrightarrow Z'$ mixing. In particular, for $z_Q = -1/3$, the coupling of the SM Higgs H to

Figure 3. Branching ratios of various two-body decay modes of Z' as functions of its mass $M_{Z'}$ for (a) $z_Q = -1/4$ and (b) $z_Q = -1/3$. In (c), we show similar BRs as functions of z_Q for $M_{Z'} = 3$ TeV. For these plots, we choose $g_z = 0.15$. Here, j includes u, d, c, s, b and ℓ includes e, μ, τ .

the Z' vanishes identically (at least at the tree level), and so does the Z - Z' mixing and, hence, these modes vanish too.¹² Less suppressed are the modes $Z' \rightarrow (H_1 + H_2)A$.

That the decays into exclusively SM modes tend to be subdominant for $|z_Q| \lesssim 1/3$ is but a reflection of the charges, which also accounts for the ratios of the said partial widths. This suppression, though, does not hold for larger values of $|z_Q|$, as is reflected by figure 3c. Indeed, for a sufficiently large $|z_Q|$, it is the dilepton or dijet modes that dominate. It is interesting to note that the value of z_Q that minimizes Z' production is not the same as when these decay modes minimize. While we could have chosen a z_Q such that, say, the dilepton signal at the LHC is minimized (rendering the model relatively free from constraints), we choose not to do so.

It is also worthwhile to consider the decay modes of $N_{1,2}$. The exact details would, of course, depend on the specific pattern of the Yukawa couplings. The decay to $W^\pm$ and $\ell^\pm$ is higher than 50% in this case because of the presence of extra U_{PMNS} factor in the decays to Z/H and ν . In calculating the decay widths, as shown in appendix, we have assumed that $M_{Z'} \gg M_N \gg M_{h^0}$.

4.2 Exclusion limits

Once the Z' has been produced at a collider, it can be detected only through its decays. The leading contribution, at the LHC, accrues from $q\bar{q}$ fusion and, for moderately large $|z_Q|$, grows as z_Q^2 . Thus, a large z_Q would facilitate detection and we deliberately choose to eschew this part of the parameter space, considering instead the case of moderately small z_Q values, when the production cross sections are not too large. Consequently, one needs

¹²This also has immediate impact on the low-energy observables, especially those measured at the Z -peak, and is reminiscent of the $B-L$ model as discussed, for example, in ref. [64].

to concentrate on decay modes that are not highly suppressed. As we have seen, for much of the parameter space of interest, $Z' \rightarrow \bar{N}_3 N_3$ is the dominant decay mode. The N_3 is not only very light, but also has a highly suppressed coupling to lighter species (the SM-like neutrinos) and, consequently, does not decay within the detector. Thus, this mode is not directly visible. However, with the emission of a visible particle (e.g., $qg \rightarrow qZ' \rightarrow q\bar{N}_3 N_3$), one could, instead have a signal comprising of a single jet accompanied by missing transverse momentum.¹³ Although the SM backgrounds to this final state is well studied, the sheer size of the background and the paucity of kinematic variables to play with renders this mode a relatively insensitive probe for a very heavy Z' .

As figure 3 shows, for such z_Q , the $\bar{N}_\alpha N_\alpha$ ($\alpha = 1, 2$) modes, together, can be competitive with the $\bar{N}_3 N_3$ mode. With the N_α decay branching fractions essentially being given by eq. (A.13), a variety of final states are possible. Particularly intriguing is the possibility of the same-sign dilepton final state, that may arise when both the N 's decay into the same sign charged lepton, with the W 's subsequently decaying into, say, jets. Also possible are the trilepton plus jets and the four-lepton final states, albeit with smaller cross sections. Many of these have been studied extensively, not only in the context of many popular new physics scenarios (such as supersymmetry or theories defined in higher dimensions), but also in those similar to ours [88–95]. A simple scaling of the cross sections convinces one that the parameter space required for reproducing the neutrino masses and mixings would be accessible only once the high-luminosity version of the LHC is operational. Rather than delve into the details thereof, we concentrate instead on the most sensitive probe.

4.2.1 From dilepton and dijet data

Despite the relatively smaller branching fraction, the decay of the Z' to a pair of charged leptons provides the strongest constraints on the parameter space, followed by the dijet signal.

Since $\Gamma(Z') \ll M_{Z'}$, expressing dilepton (or dijet) production in terms of an on-shell Z' production followed by its decay constitutes an excellent approximation. The leading order contribution emanates from $q\bar{q}$ fusion, and has a simple structure as given in eq. (4.1). We, though, include the next-to-leading order QCD corrections, parametrizable in terms of a K -factor of 1.3 [96]. As for the parton fluxes, we use the NNPDF2.3LO [97] parton distributions, with the natural choice for the renormalization and factorizations scales, namely $\mu_F = \mu_R = M_{Z'}$.

Exclusion bounds on the model parameters, from a given experiment, can be obtained by comparing the expected signal strength with the upper bound (UB) on new physics events that the non-observation of an excess in the said experiment implies. To this end, we use the dilepton [98, 99] and dijet [100, 101] resonance search data that the two LHC experiments have collected (at $\sqrt{s} = 13$ TeV) with an approximate integrated luminosity of 140 fb^{-1} . We start by summarizing the experimental results:

- **ATLAS dilepton [98]:** The ATLAS collaboration has performed a high-mass spin-1 resonance search in the dilepton final state in the mass range of 0.25 TeV to 6 TeV

¹³Similarly, monophoton, mono- Z or mono- W signals (accompanied, in each instance by a transverse momentum imbalance) are possible too, but these suffer from additional coupling constant suppressions.

with an integrated luminosity of 139 fb^{-1} . We recast their upper bound on fiducial $\sigma \times BR$ for a spin-1 selection with a width/mass hypothesis of 1.2% as is applicable for our analysis.¹⁴ The definition of the fiducial phase-space region and the fiducial selection efficiency can be found in ref. [98]. While this efficiency differs slightly for the dielectron and dimuon channels and varies, in addition, with the resonance mass, for the sake of simplicity, we use a fixed fiducial selection efficiency of 0.6 for the entire dilepton invariant mass range.¹⁵ We obtain the observed $\sigma \times BR$ UB from the HEPDATA [102] repository.

- **ATLAS dijet** [100]: For the dijet channel (also done with the same luminosity), the collaboration presents an UB on $\sigma \times BR \times A$ (where the acceptance A can be approximated to 0.4). Recasting the data presented in figure 8a of [100] for a generic Gaussian signal in the inclusive channel with a 3% width/mass hypothesis, we obtain our exclusion limits.
- **CMS dilepton** [99]: The CMS collaboration has performed a high-mass spin-1 resonance search in the dilepton final state in the mass range from 0.2 TeV to 5.4 TeV. Working with an integrated luminosity $\sim 140 \text{ fb}^{-1}$, they present an upper bound on $\sigma \times BR$, assuming the SM value for the width/mass ratio, namely 0.6%.
- **CMS dijet** [101]: For the high-mass dijet events, the collaboration uses an integrated luminosity of 137 fb^{-1} . We recast the observed UB on $\sigma \times BR \times A$ (with $A = 0.5$) taken from figure 10 (lower panel) of ref. [101] for the spin-1 resonance with a width/mass ratio of 1%. In all these four searches discussed above, we use the ones with smallest width/mass ratio which are available in those analyses as in our case Z' width is much smaller compared to its mass.

As a particular example, we display, in figure 4, the exclusion limits on the Z' mass, as gleaned from the ATLAS dilepton resonance search data, for two choices of z_Q , namely $z_Q = -1/4$ and $-1/3$. Working with the fiducial $\sigma_{fid} \times BR$ as provided in the ATLAS paper [98], we obtain $\sigma \times BR$ using a fiducial selection efficiency of 0.6. The lower limits on $M_{Z'}$, for $g_z = 0.1$, are about 2.2 (2.6) TeV for $z_Q = -1/4$ ($-1/3$). The difference in the two exclusions can be traced to two factors, namely a slightly smaller Z' -production cross section, as well as a slightly smaller branching fraction into a charged lepton pair. Note that (as promised earlier) the dependence on z_Q is not too severe. Were one to be interested in the mass exclusions for other z_Q and g_z values, these could be obtained trivially by realizing that the production cross section scales as g_z^2 , and reading off the dilepton BR from figure 3.

¹⁴Ref. [98] has performed the analysis for several values of this ratio, and we choose the one closest to our situation. We have checked that our conclusions are not too sensitive to this choice.

¹⁵We have checked that the consequences of this approximation are too small to be relevant.

Figure 4. Comparison of the 95% CL upper bound on the observed and the expected $\sigma(pp \rightarrow Z') \times BR(Z' \rightarrow \ell\ell)$ obtained from the ATLAS dilepton resonance search data at the 13 TeV LHC with $L = 139 \text{ fb}^{-1}$ with the theoretical predictions of our model for $z_Q = -1/4$ and $-1/3$ choices. We use the reference value for the $U(1)_z$ gauge coupling $g_z = 0.15$. The green and yellow bands represent the 1σ and 2σ uncertainty regions of the expected values respectively.

4.2.2 Low-energy observables

A nonzero value of z_H , the $U(1)_z$ charge of the SM Higgs doublet, induces tree-level $Z \leftrightarrow Z'$ mixing. This has two main ramifications:

- Tree-level contributions to the oblique parameters are induced. In particular, the tree-level contribution to the T -parameter is given by [80],

$$\alpha_{EM} T^{\text{new}} = \frac{\Pi_{ZZ}^{\text{new}}}{M_Z^2} = \frac{M_Z^2 - (M_Z^0)^2}{M_Z^2}. \quad (4.4)$$

Here, M_Z is the Z -boson mass in the new theory and, for our purposes, it suffices to consider the tree-level expression as given in eq. (2.7). Similarly, M_Z^0 , the mass within the SM, is given by $M_Z^0 = g_w v_h / 2 \cos w$ at the tree-level. And, finally, α_{EM} denotes the fine-structure constant at Z -pole. In effecting the actual calculation, the higher-order SM contributions would, of course, have to be taken into account, and we have done so. On the other hand, loop corrections to T wrought by new physics are further suppressed by large masses and the $Z \leftrightarrow Z'$ mixing angle and can be safely neglected. We use the value $T = 0.07 \pm 0.12$ ref. [26] in our analysis.

- A related constraint arises from the measurement of the Z -coupling of the light fermions, occasioned, again, primarily by the $Z \leftrightarrow Z'$ mixing. Determined from the

forward-backward asymmetry or through the line-shape of the Z -resonance, these observables as also the Z -width [80] are very precisely measured [26].

- Another relevant constraint can come from the LEP measurements. The Z' boson, despite being heavier than the LEP energies, can contribute to the $e^+e^- \rightarrow \bar{f}f$ processes through the interference with the γ and Z mediated processes. For the sequential-SM, the 95% confidence level lower limit on the Z' mass is 1760 GeV as obtained from the LEP data [103]. In our case, this limit is much more relaxed to the point of being irrelevant since the Z' couplings with leptons and quarks are much smaller than the SM-like couplings for our benchmark parameters. Therefore, we do not consider LEP constraints in our analysis.

Once again, due to the vanishing of z_H for $z_Q = -1/3$, there is no tree-level $Z \leftrightarrow Z'$ mixing, and neither of the aforementioned constraints are applicable, at least at the tree-level. The one-loop effect is too small to be of any consequence. As for $z_Q = -1/4$, our choice commensurate with neutrino phenomenology as well as possible unification, owing to it not being too far from $-1/3$, the $Z \leftrightarrow Z'$ mixing is still very small, and such low-energy observables do not strongly constrain the parameter space.

We show the exclusion plots in the $M_{Z'} - g_z$ plane in figures 5a and 5b for $z_Q = -1/4$ and $z_Q = -1/3$ choices respectively, while in figure 5c, we show the similar exclusion in the $M_{Z'} - z_Q$ plane for fixed $g_z = 0.15$. In each case, we show exclusion regions using the latest dilepton, dijet resonance search data from the LHC, as well as those coming from the T -parameter and Z -width measurements. As expected, the dilepton data does impose severe constraints. However, owing to $Br(Z' \rightarrow \bar{\ell}\ell)$ assuming its minimum at $z_Q \sim -1/4$, the lower limit on $M_{Z'}$ for both values of z_Q considered here is as low as ≈ 2 TeV for $g_z \sim 0.15$. And while the dijet branching fraction is comparable to the dilepton one, this mode suffers from a much larger (QCD) background, and consequently the bounds are much weaker. As for T -parameter and Z -width measurements, since tree-level $Z \leftrightarrow Z'$ mixing is absent for $z_Q = -1/3$, the constraints are virtually nonexistent, and continue to be very weak for $z_Q = -1/4$ as well.

5 A dark matter candidate?

As we have already seen in section 2 and section 3, the model may have three ultralight particles, *viz.*

- a pseudoscalar pseudo-Nambu-Goldstone boson A ,
- a largely singlet neutrino N_{3R} , and
- an even lighter doublet-like neutrino ν .

The lightest of the three, being stable, would be an apparent candidate for the dark matter. Indeed, given the lifetimes and interaction strengths, the “dark sector” could be much richer.

Figure 5. Exclusion regions in the $M_{Z'} - g_z$ plane for fixed (a) $z_Q = -1/4$ and (b) $z_Q = -1/3$ and in (c) in the $M_{Z'} - z_Q$ plane for fixed $g_z = 0.15$. We show exclusion regions using T -parameter, Z -width, and the latest dilepton and dijet data from the LHC.

It is, of course, well-known that a light doublet-like neutrino would have decoupled while still relativistic and would constitute hot dark matter. Since this would have interfered with large-scale structure formation, a hot dark matter can constitute only a minor fraction of the total dark matter relic density and, fortunately, the very structure of the SM ensures that the usual neutrinos satisfy this condition. And with ν being overwhelmingly doublet-like, it too would automatically satisfy the same.

This allows us to direct our attention to the pseudoscalar A and the singlet-like neutrino N_{3R} . As we have already discussed, the former's mass can be trivially uplifted by the dint of adding a third singlet χ_3 such that nontrivial trilinear and/or quadrilinear terms are admissible in the scalar potential. While either of soft and hard breaking of the global symmetry — down to a single $U(1)$ which is gauged — will render the pseudoscalar(s) massive, in the case of the former (soft breaking), the resultant mass is controllable (and we have a pNGB). This constitutes a particularly simple strategy as the neutrino masses and mixings are essentially left unchanged. As for collider signals, while the partial widths $Z' \rightarrow A\xi_i$ would be altered on account of the A picking up a mass, the changes are not very significant for relatively small M_A values. More importantly though, the A may now decay within the detector, thereby eliminating this particular source of a missing transverse energy signal. This, however, would be replaced by more exotic (and, hence, more visible) signals at the LHC, perhaps including displaced vertices.

Concentrating on the N_{3R} (-like) state, let us begin by examining its mass. As a perusal of eq. (3.4) shows, the largest contribution to this mass would accrue from the seesaw like mechanism involving the N_{1R} and N_{2R} fields. This would lead to an effective mass for the mass-eigenstate Ψ of¹⁶

$$m(\Psi) \sim \left(s_{\alpha 3} \frac{x_1^4 x_2}{\Lambda^4} \right)^2 \frac{\Lambda}{x_2^2} \sim s_{\alpha 3}^2 \xi^7 x.$$

For $s_{\alpha 3}$ of around 0.05 and the two heavy RHNs of the order 1.2 TeV, we get the mass of N_3 to be a few keVs. A sterile neutrino at such a scale is particularly interesting from the dark matter perspective because it can have a very long lifetime, comparable to the age of the Universe. For the keV scale DM candidates, almost the entire observed relic density can be accounted for only if $m_{\text{DM}} > 0.4$ keV [104, 105], which seems plausible in our case. Still, large mixings with neutrinos, while being a desirable feature from the collider search point of view, can end up producing way more dark matter than we require, thereby potentially spoiling its DM candidature.

A state as light as this can only have two types of decays, namely $\Psi \rightarrow 3\nu$ (where ν are the SM-like mass eigenstates) and $\Psi \rightarrow \nu_i + \gamma$. Clearly, the latter are loop-suppressed and, hence, of relatively little significance. As for the former set of modes, it is instructive to consider, formally, the lowest-dimensional operator within the effective theory that can lead to them. This is most easily done in terms of the gauge eigenstates, *viz.*,

$$\frac{a_1}{\Lambda^4} (\bar{\ell}_i \gamma_\mu \ell_j) (\bar{\ell}_k \gamma^\mu N_{3R}^c) H^* \chi_2.$$

For $a_1 \sim \mathcal{O}(1)$, this leads to

$$\tau_\Psi \gtrsim \frac{10^{25}}{a_1^2} \text{ s} \times \left(\frac{100 \text{ MeV}}{m_\Psi} \right)^5,$$

which is comfortably larger than the age of the universe ($\tau_U \sim 5 \times 10^{17}$ s). Naively, it might seem then that we might have as well admitted a different structure that would have allowed for a much heavier N_{3R} .

¹⁶Note that $m(\Psi)$ is still much larger than the s_{33} term in eq. (3.4), a situation very analogous to that in the doublet-sector, namely the large difference between $M_{3 \times 3}$ and the corresponding Weinberg operator.

This, however, is misleading. As we have seen in section 3, the diagonalization of the mass matrix leads to mixings between the ν -like and the N_3 like eigenstates. In other words, the mass eigenstates Ψ, ν' are symbolically given by

$$\Psi \approx \cos \theta_i N_3 + \sin \theta_i \nu_i, \quad \nu' \approx -\sin \theta_i N_3 + \cos \theta_i \nu,$$

where generational dependence has been omitted. Other mixings connecting ν_i, N_{3R} with $N_{1,2}$ are not taken into account as they are much suppressed and, therefore, not significant for the issue at hand.

Owing to the fact that the N_{3R} is a singlet, this mixing immediately leads to a $Z\bar{\Psi}\nu'$ coupling which, in turn, leads to a Z -mediated contribution to the $\Psi \rightarrow 3\nu$ decay. The corresponding partial widths are

$$\Gamma_i = \Gamma(\Psi \rightarrow \nu_i \bar{\nu}_j \nu_j) \sim \frac{G_F^2 M_{N_3}^5}{192\pi^3} \sin^2 \theta_i \left(1 - \frac{\delta_{ij}}{2}\right), \quad (5.1)$$

leading to a lifetime $\tau_\Psi \gtrsim 10^{18}\text{s}$, for $m_N \sim 6\text{keV}$ and mixing of $\mathcal{O}(10^{-2})$. So for such masses and mixings, we still can barely have $\tau_\Psi > \tau_U$. This ostensible enhancement compared to the earlier estimate can be traced back to the fact that several powers of (x/Λ) in the aforementioned effective operator are actually subsumed in the suppression of m_{N_3} itself, and must not be double counted.

At this point, several issues need to be delved upon. While it might seem that we naturally have $\tau_\Psi > \tau_U$, in reality, τ_Ψ depends crucially on the values of the Wilson coefficients, and, thus, such a requirement on the lifetime imposes conditions on the WCs. It should be appreciated though that, for such a light DM particle, $\tau_\Psi > \tau_U$ is not a strict requirement. In fact, a τ_Ψ value even somewhat smaller than τ_U may also be admissible for this would only mean that a fraction of the DM has decayed in the course of the Universe's evolution yet leaving behind sufficient relic density. Since the decay is into neutrinos alone, the only discernible effect would be through altering the neutrino-photon ratio in the early universe, thereby altering the effective number of relativistic degrees of freedom, and thus invite constraints from this measurement. However, the existence of such a restriction depends crucially not only on the epoch of this decay, but also on specific mechanism (e.g., freeze-in versus freeze-out) of DM relic density generation, and is not immediately applicable to the case at hand.

On the other hand, even if a τ_Ψ value somewhat smaller than τ_U is obtained, these same operators would lead, at one loop, to $\Psi \rightarrow \nu_i + \gamma$ manifesting itself through X-ray lines. The non-observation of such signal in different low energy experiments, mandates [106, 107] that the sterile-SM neutrino mixing angle satisfies $\theta_i \ll 10^{-2}$ if a 6 keV Ψ were to provide full DM relic density. The simplest way to satisfy such seemingly incompatible constraints would then be to assume that the Ψ provides for only a small fraction of the relic density in the form of a potentially warm component, especially since for certain regions of the parameter space, it could be produced non-thermally. More interestingly, the Ψ has considerable self-interaction (mediated by the Z') with a suppressed, but long-distance component mediated by the pseudoscalar A . This has the potential to provide some pressure to the DM fluid and, thereby, address certain long-standing issues pertaining to details of structure formation.

To anoint Ψ to be the main or even a significant DM constituent, one must ascertain not only whether the correct relic density can be reached but also whether the scenario falls foul of other constraints, both cosmological as well as those emanating from laboratory tests (both direct and indirect detection). This demands detailed analysis that is beyond the scope of the present work. However, at the same time, we want to emphasize a few general issues. In our quantitative analysis we made several simplifying assumptions regarding x_i , Λ and some of the WCs. Tweaking these assumptions can substantially change the masses and mixings of the neutrino, as is seen in the context of a Froggatt-Nielsen scenario through the introduction of different scales through different powers of a scaling factor [108]. In a similar vein, by altering the $U(1)_z$ charges of the N_i (while maintaining anomaly cancellation) and/or the scalar fields χ_a , the neutrino mass matrix can be changed. This would allow us a much larger mass for the N_3 -like state, *viz.* $\mathcal{O}(\text{MeV})$, with further suppressed N_3 - ν_i mixings thereby still allowing for $\tau_\Psi \gtrsim \tau_U$. Consequently, the standard freeze-in mechanism would hold for such a DM. We would like to postpone these issues to a future project.

To examine the falsifiability of our hypothesis, it is important to consider the strength of the interactions that mediate low-energy scattering involving Ψ and the SM particles. Fortunately, such interactions are not unduly suppressed thanks to the fact that both N_{3R} and the SM fermion carry $U(1)_z$ quantum numbers. Consequently, the interaction strength is governed only by $g_z^2/M_{Z'}^2$, or, equivalently, by x_i^{-2} , and in the case of freeze in DM generation, relic density is proportional to this interaction, that, $M_{Z'}$ being at the TeV-scale, exactly represents an example of how entire relic density can be reproduced. Rather than present a full analysis, we refer the reader to the existing literature. For example, it has been shown, in an analogous context, in ref. [109] that a parameter space, consistent simultaneously with the requisite relic density, the measurement of the cosmologically relevant effective relativistic degrees of freedom and energy injection, from DM annihilation, into the cosmic microwave background radiation, can be found. The required interaction strength is of the same order as what transpires naturally in our model and would leave such particles undetectable in the currently operative (satellite-based) indirect detection experiments [109]. Even more interestingly, such a DM is likely to be detectable not only at the next generation of direct detection experiments, but also at the Super-Belle detector [110].

Before we end this section, we would like to remind the reader of a possibility that we did not elaborate on. Consider, for example, the case where the extra global $U(1)$ is not broken by terms in the potential. The Yukawa couplings, nonetheless, do break it and, consequently, quantum corrections¹⁷ would lift the mass of the Goldstone, by an amount dictated by the effective Yukawa couplings. A naive perusal of the eq. (3.4) suggests that the largest such corrections to the mass of CP-even component of the χ_i is $\delta m \sim \mathcal{O}(s_{\alpha\beta}x_2/4\pi)$. This, however, is not the correction that accrues to the mass of the pNGB. The latter can only be described in terms of a Coleman-Weinberg-like potential that breaks this accidental $U(1)$ symmetry, and is, understandably, orders of magnitude smaller. For $m_A > 2m_\Psi$, the

¹⁷As is obvious, and has been said in section 2.2, and in particular in footnote 5, the gauge boson or scalar loops would have vanishing contributions owing to the fact that this extra $U(1)$ is exact as far as those sectors are concerned.

pNGB can decay into the DM, while for $m_A < 2m_\Psi$, it is very long-lived. Free from restrictions (such as those imposed by X-ray or Lyman- α observations), this could, again, play a significant role in the evolution of galaxy clusters etc. [111]. A detailed examination of such effects is beyond the scope of this paper and is postponed for a future study.

6 Summary and conclusion

With the aim of explaining neutrino masses without invoking either ultrasmall Yukawa couplings or an almost inaccessible new (seesaw) scale, we consider a scenario where the gauge symmetry has been augmented by an extra $U(1)_z$. If its action on the SM particles is nontrivial, but generation-invariant (so as to allow for a single SM Higgs to give masses to the charged leptons), then the possible charge assignment for the right-handed neutrinos (RHN) is severely restricted by the requirement of gauge (and mixed gauge-gravity) anomaly cancellation. (We assume here that, unlike in certain popular schemes such as the inverse seesaw mechanism, we have the minimum possible number of RHNs.) Only the most trivial such assignment allows for tree-level neutrino Dirac mass terms. On the other hand, bare Majorana mass terms cannot be incorporated. Indeed, analogues of the Weinberg term can be written only if the new Higgs breaking the $U(1)_z$ have one of two specific choices of the charge.

For any choice of the RHN charges other than the most trivial one, not only are renormalizable Dirac mass terms disallowed, but so are the Majorana mass terms except for specific choices of the $U(1)_z$ breaking Higgs bosons. Completely unrelated to this, the absence of any resonance in the LHC data has pushed the mass of the new gauge boson Z' to above several TeVs.

In view of this, we assume an agnostic standpoint claiming that any such theory can, at best, be the low-energy limit of a more fundamental theory, characterised by a cut-off scale Λ . This, immediately, allows us to write non-renormalizable terms suppressed by powers of Λ . While a wide variety of such terms, in principle, can be written, we concern ourselves only with the neutrino sector. Invoking the next to the trivial quantum number assignment for the RHNs, we then write down all relevant higher-dimensional terms *à la* the Froggatt-Nielsen mechanism. Using the power of higher dimensional operators to the hilt, we generate tiny neutrino masses without any need to invoke tiny Yukawa couplings. Indeed, even without using all the free parameters of the theory, it can naturally reproduce the experimentally observed neutrino mixings and mass-squared differences, while satisfying the cosmological bound on the sum of masses as well as that from non-observation of neutrinoless double-beta-decay. Simultaneously, it prophesies, amongst others

- a pair of heavy RHNs $N_{1,2}$ at the 1 TeV mass scale that decay promptly into ℓW , νZ and promise interesting signals at the high-luminosity run of the LHC;
- a moderately heavy Z' ($m_{Z'} \lesssim 3 \text{ TeV}$ for $g_{Z'} \sim 0.15$) that escapes LHC bounds — from dilepton and dijet searches — despite having unsuppressed couplings with the quarks and leptons, simply by virtue of decaying primarily into the RHNs. Similarly, for natural choices of $U(1)_z$ charges (especially those commensurate with possible

charge quantization), the LEP constraints such as those on the oblique parameters are trivially satisfied;

- a light RHN N_3 in the keV — MeV range. With the Z' having a large branching fraction into a N_3 -pair, and with the N_3 being stable at the collider timescales, this would lead to additional contribution to the monojet (monophoton) plus missing transverse momentum signal at the LHC; Indeed, for a large part of the parameter space, the N_3 can have a lifetime comparable to or even greater than that of the Universe and, thus, can constitute a warm DM component.
- a pseudoscalar pseudo-Nambu Goldstone boson, with its mass uplifted only by quantum corrections or additional soft terms in the scalar potential (the latter being absent in the simplest realization). This has the potential of being an additional contributor to the DM relic density (while escaping many of the constraints applicable to N_3). Furthermore not only does it have non-negligible self-interaction, but it can also mediate N_3 scattering thereby playing an important role in not only determining the relic density, but also in engendering a non-negligible pressure term for the DM fluid and thereby affecting the details of structure formation.

The model presented, thus, offers much more than an understanding neutrino phenomenology. Not only does it offer tantalizing prospects at the LHC, but also intriguing avenues to explore in the context of dark matter and details of structure formation. We hope to return to more in-depth study of these issues in a future publication.

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A Decay widths

A.1 Decay widths of Z'

In this appendix, we provide the analytical expressions of the tree-level partial widths of various two-body decay modes of Z' . These expressions are computed using the Feynman rules obtained from the interaction Lagrangian shown before them.

- $Z' \rightarrow \bar{f}f$: For the following interaction Lagrangian,

$$\mathcal{L}_{Z'\bar{f}f} = g_L \bar{f}_L \gamma^\mu f_L Z'_\mu + g_R \bar{f}_R \gamma^\mu f_R Z'_\mu, \quad (\text{A.1})$$

the expression for the $Z' \rightarrow \bar{f}f$ partial width is given by

$$\Gamma_{(Z' \rightarrow \bar{f}f)} = \frac{N_c M_{Z'}}{24\pi} \sqrt{1 - \frac{4M_f^2}{M_{Z'}^2}} \left[(g_L^2 + g_R^2) \left(1 - \frac{M_f^2}{M_{Z'}^2} \right) + 6g_L g_R \frac{M_f^2}{M_{Z'}^2} \right]. \quad (\text{A.2})$$

In the above, g_L and g_R are the left- and the right-handed couplings respectively, M_f is the mass of the fermion f and N_c is the corresponding number of colors.

- $Z' \rightarrow \nu_R \nu_R$: For the following interaction Lagrangian,

$$\mathcal{L}_{Z' \nu_R \nu_R} = g_\nu (\nu_R^c)^T \gamma^\mu \nu_R Z'_\mu, \quad (\text{A.3})$$

the expression for the $Z' \rightarrow \nu_R \nu_R$ partial width is given by

$$\Gamma_{(Z' \rightarrow \nu_R \nu_R)} = \frac{M_{Z'}}{24\pi} g_\nu^2 \left(1 - \frac{4M_{\nu_R}^2}{M_{Z'}^2} \right)^{3/2}, \quad (\text{A.4})$$

where g_ν is the coupling and M_{ν_R} is the mass of the RHN.

- $Z' \rightarrow W^+ W^-$: For the following triple gauge boson interaction with strength λ_W ,

$$\mathcal{L}_{Z' W^+ W^-} \supset \lambda_W Z'_\mu (p_1) W_\nu^+(p_2) W_\rho^-(p_3), \quad (\text{A.5})$$

the expression for the $Z' \rightarrow W^+ W^-$ partial width is given by

$$\Gamma_{(Z' \rightarrow W^+ W^-)} = \frac{M_{Z'}^5}{192\pi M_W^4} \lambda_W^2 \left(1 - \frac{4M_W^2}{M_{Z'}^2} \right)^{3/2} \left(1 + \frac{20M_W^2}{M_{Z'}^2} + \frac{12M_W^4}{M_{Z'}^4} \right). \quad (\text{A.6})$$

- $Z' \rightarrow ZS$: From the following interaction with dimensionful coupling strength μ_S ,

$$\mathcal{L}_{Z' Z S} = \mu_S Z'_\mu Z^\mu S, \quad (\text{A.7})$$

where S is a CP even scalar, the expression for the corresponding partial width is given by

$$\begin{aligned} \Gamma_{(Z' \rightarrow ZS)} &= \frac{\mu_S^2 M_{Z'}}{192\pi M_Z^4} \left(1 - \frac{(2M_S^2 - 10M_Z^2)}{M_{Z'}^2} + \frac{(M_S^2 - M_Z^2)^2}{M_{Z'}^4} \right) \\ &\times \left(1 - \frac{2(M_S^2 + M_Z^2)}{M_{Z'}^2} + \frac{(M_S^2 - M_Z^2)^2}{M_{Z'}^4} \right). \end{aligned} \quad (\text{A.8})$$

- $Z' \rightarrow SA$: From the following interaction with cubic coupling

$$\mathcal{L}_{Z' S A} = g_p Z'_\mu \partial^\mu S A, \quad (\text{A.9})$$

where S is a CP -even scalar and A is a CP -odd scalar

$$\begin{aligned} \Gamma_{(Z' \rightarrow ZS)} &= \frac{g_p^2}{12\pi M_{Z'}^5} \left[(M_S - M_A - M_{Z'}) (M_S + M_A - M_{Z'}) \right. \\ &\quad \left. (M_S - M_A + M_{Z'}) (M_S + M_A + M_{Z'}) \right]^{3/2} \end{aligned} \quad (\text{A.10})$$

A.2 Decay widths of heavy RHN

The heavy RHN decay modes are given by the Lagrangian of the form:

$$\mathcal{L} = -\frac{g_w}{\sqrt{2}}\bar{l}_L\gamma_\mu U_{\nu N}NW^\mu - \frac{g_w}{2\cos w}\bar{\nu}\gamma_\mu U_{\nu\nu}^\dagger U_{\nu N}NZ^\mu - \frac{H}{v_h}M_N^{\text{diag}}U_{\nu N}^\dagger U_{\nu\nu}\nu + \text{h.c.} \quad (\text{A.11})$$

The decay rates are then given by:

$$\Gamma(N_\alpha \rightarrow W^- l_i^+) = \Gamma(N_\alpha \rightarrow W^+ l_i^-) \approx \frac{g_w^2}{64\pi M_W^2} M_N^3 |(U_{\nu N})_{i\alpha}|^2, \quad (\text{A.12})$$

$$\Gamma(N_\alpha \rightarrow Z\nu_i) \approx \Gamma(N_\alpha \rightarrow H\nu_i) \approx \frac{g_w^2}{64\pi M_W^2} M_N^3 |(U_{\nu\nu}^\dagger U_{\nu N})_{i\alpha}|^2, \quad (\text{A.13})$$

where $U_{\nu\nu}$ is approximately the *PMNS* matrix. $U_{\nu N}$ is the mixing between the light SM neutrinos and the heavy RHNs given by [112, 113]

$$U_{\nu N} = -\mathcal{D}(M_N^{\text{diag}})^{-1}. \quad (\text{A.14})$$

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