Motivated by the ongoing searches for new physics at the LHC, we explore the low energy consequences of a D-brane inspired $SU(4)_C \times SU(2)_L \times SU(2)_R$ (4-2-2) model. The Higgs sector consists of an $SU(4)$ adjoint, a pair $H + \bar{H}$ in $\begin{pmatrix} 4 & 1 & 2 \\ 1 & 2 & 2 \end{pmatrix}$, and a bidoublet field in $h(1, 2, 2)$. With the $SU(4)$ adjoint the symmetry breaks to a left-right symmetric $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ model. A missing partner mechanism protects the $SU(2)_R$ Higgs doublets in $H, \bar{H}$, which subsequently break the symmetry to the Standard Model at a few TeV scale. An inverse seesaw mechanism generates masses for the observed neutrinos and also yields a sterile neutrino which can play the role of dark matter if its mass lies in the keV range. Other phenomenological implications including proton decay are briefly discussed.
1 Introduction

Spontaneously broken left-right symmetry and the idea of quark-lepton unification are nicely encapsulated in the $SU(4)_C \times SU(2)_L \times SU(2)_R$ (4-2-2) gauge symmetry proposed by Pati and Salam [1]. With gauge bosons that can mediate proton decay absent, the 4-2-2 symmetry may well be broken at scales that are orders of magnitude lower than the usual grand unified scale $M_{GUT} \simeq 2 \times 10^{16}$ GeV. A lower bound on the 4-2-2 breaking scale of around 100 TeV arises in the simplest model from the non-observation of the decay $K \rightarrow \mu + e$ mediated by some of the gauge bosons in $SU(4)_C$. The experimental constraints on the left-right gauge symmetry breaking scale are significantly milder [2, 3], and indeed this symmetry may be broken at a scale that may be accessible either at the LHC or its upgrades.

In this letter we propose to explore the low energy consequences of a 4-2-2 model derived from an intersecting D-brane framework. As is well known, the 4-2-2 model realises quark-lepton unification in a natural way. Several other attributes of this symmetry [4]-[7] are the incorporation of the right-handed neutrino in the spectrum, the absence of gauge bosons mediating fast proton decay etc. The last few years there has been growing theoretical and experimental interest in the theory of fundamental interactions in the TeV region. We believe that the 4-2-2 model incorporates all the ingredients to interpret possible related findings and therefore in this note, we would like to examine low scale symmetry breaking patterns and explore possible predictions. Indeed, the present experimental bounds on scalar superpartners are close to the TeV scale while searches for new gauge bosons have put lower limits at a few TeV. Supersymmetric scalar masses of the MSSM states are definitely related to supersymmetry breaking, while the existence of new gauge bosons are naturally associated with some new symmetry breaking scale. A natural candidate of such a new gauge symmetry, not far above the electroweak (EW) scale, is $SU(2)_R$, so that above this breaking scale, the model is left-right symmetric. Moreover, there is a good chance that the associated gauge boson with a relatively low $SU(2)_R$ breaking scale leaves its signature in experiments through new interactions, which may be identified in future Diboson searches [9]. The left-right symmetric model is naturally embedded in a quark-lepton unified (4-2-2) symmetry which admits an interesting string realisation in the context of heterotic superstrings, the intersecting D-brane scenarios, as well as in F-theory [10]-[13]. However, although the PS symmetry unifies quarks and leptons, yet its symmetry structure consists of a product of non-abelian factors and as such, does not necessarily imply unification of the gauge couplings. Given this fact, from the above string theory realisations, the intersecting D-brane scenarios, as well as in F-theory [10]-[13]. However, although the PS symmetry unifies quarks and leptons, yet its symmetry structure consists of a product of non-abelian factors and as such, does not necessarily imply unification of the gauge couplings. Given this fact, from the above string theory realisations, the intersecting D-brane set up, appears to be the most natural scenario. Indeed, within this context, each group factor is associated with a different brane stack with its own gauge coupling strength, therefore the PS model is naturally motivated in an intersecting D-brane framework. Moreover, the tight connection between the unification scale and the Planck scale $M_{Planck}$ is no longer mandatory and the symmetry breaking scale can be substantially lower than $M_{Planck}$.

In this note, we investigate the realisation of a two step symmetry breaking scenario of the

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[^1]: For a recent review see [8].
PS model built in the context of open string theories and D-branes. At the first stage the PS symmetry breaks at $M_{GUT}$ to the left-right (LR) symmetric model with the use of the adjoint Higgs. At a second scale $M_R \ll M_{GUT}$ a pair of right-handed Higgs doublets triggers the breaking of the LR-symmetry down to the Standard Model one.

The layout of the paper is as follows. In section 2 we discuss the D-brane framework for the 4-2-2 model and present the particle spectrum in Table 1. A variety of new fields appear that are absent in the standard field theory constructions. They include a SM sterile neutrino which turns out to be a plausible dark matter candidate. In section 3 we consider the two step breaking of 4-2-2 to the SM. We exploit here a \textit{triplet-singlet} splitting mechanism, which is analogous to the well-known \textit{doublet-triplet} mechanism of SU(5). The left-right symmetry breaking scale lies in the TeV range which may be found at the LHC or its future upgrades. Section 4 contains a discussion of neutrino masses which includes the inverse seesaw mechanism and a sterile keV mass neutrino. A brief discussion in this section of proton decay shows that it is adequately suppressed in this class of models. Our summary and conclusions are presented in section 5.

2 4-2-2 Spectrum from intersecting D-branes

In this section we will present the basic features of an open string realisation of the 4-2-2 gauge symmetry. In general, there are several methods to construct 4-2-2 vacua in string theory, such as Calabi-Yau compactification, Type IIA string theory, interacting CFT constructions and also Gepner constructions. For the purposes of this work we find it convenient to represent these models in terms of intersecting D-brane configurations which is the appropriate description for type IIA string theory.

The low energy phenomenology of the 4-2-2 symmetry built in the framework of intersecting D-brane scenarios differs in many respects from the corresponding model which admits an SO(10) embedding \cite{7}. The main reason is that intersecting D-brane constructions yield an $U(4) \times U(2) \times U(2)$ gauge group, which is the 4-2-2 symmetry augmented by three $U(1)$ factors. Indeed, recalling that $U(n) \simeq SU(n) \times U(1)/Z_n$ the final gauge symmetry of the D-brane version of the effective field theory model is

$$SU(4)_C \times SU(2)_L \times SU(2)_R \times U(1)_C \times U(1)_L \times U(1)_R.$$ (1)

The intersecting D-brane set up associated with the above symmetry is depicted in Figure 1. Each gauge group factor $U(n)$ is associated with a set of $n$ parallel, almost coincident, D-branes, while the various massless states are represented by open strings attached on the various sets of branes in the appropriate configurations. This set up gives rise to bi-fundamental representations which accommodate the particle spectrum of the model. The most general picture of the available representations is as follows: i) Open string connecting two brane stacks, give rise to bi-fundamental representations with respect to the corresponding gauge groups; ii) open strings stretched between a D-brane and its image transform in the antisymmetric or symmetric representation of the gauge group; ii) for each group factor, we should include the adjoint
representation which arises from open strings stretched between branes in the same stack.

To underline the salient features of the D-brane derived 4-2-2 models, in Table 1 we present a minimal set of fields obtained in the context of a D-brane set up, which are required for the realisation of a viable effective model. There are states originating from string vibrating at the intersections \(ab, bc, ca\) of the three D-brane stacks \(a, b, c\). There are also states emerging from strings with ends attached to the mirror branes \(a^*, b^*, c^*\) (not shown in figure). The observed chirality of the spectrum can be adjusted from appropriate string boundary conditions but, in general, the spectrum may contain additional vector like pairs.

![Figure 1: Intersecting D-brane configuration for the 4-2-2 symmetry. The letters \(a, b, c\) respectively stand for the \(U(4)_C, U(2)_L\) and \(U(2)_R\) brane stacks. Under orientifold planes there are the mirror branes (not shown in this figure) denoted with \(a^*, b^*, c^*\) in the main text.](image)

As expected, all massless states associated with the open string spectrum, are charged under the extra \(U(1)\) factors. Indeed, returning to the content of Table 1 we note that the last three entries show the ‘charges’ of the representations under the three abelian factors \(U(1)_C, U(1)_L, U(1)_R\). It is remarkable that the abelian factor \(U(1)_C\) plays a central rôle in defining the baryon and lepton quantum numbers. Indeed, if we consider the breaking

\[
U(4)_C \to SU(4)_C \times U(1)_C \to SU(3) \times U(1)_{B-L} \times U(1)_C ,
\]

we find out that they are defined in terms of the linear combinations

\[
U(1)_B = \frac{1}{4} (U(1)_C + U(1)_{B-L}) ,
\]

\[
U(1)_L = \frac{1}{4} (U(1)_C - 3U(1)_{B-L}) .
\]

We note that in the context of the Standard Model these are global symmetries which are anomalous. In the framework of intersecting D-brane constructions, the \(U(1)_{C,L,R}\) symmetries also are anomalous and only particular combinations are anomaly-free. It transpires that these \(U(1)\) anomalies are cancelled by the generalised Green-Schwarz mechanism (through couplings
between the RR two form field and the corresponding field strength, see for example \[13\]), which induces masses for the corresponding gauge bosons. Yet, at the perturbative level, these remain as global symmetries which prevent rapid baryon and lepton number violating processes.

In the following we describe the embedding of the Standard Model states in the 4-2-2 representations and the Higgs mechanism for a two-stage symmetry breaking to SM. The fermion generations are accommodated in the following representations

\[
F_L + \bar{F}_R = (4, 2, 1) + (\bar{4}, 1, 2),
\]

which make up just the 16 of the $SO(10)$. The Standard Model particle assignment is

\[
F_L = (4, 2, 1) = Q(3, 2, \frac{1}{6}) + \ell(1, 2, -\frac{1}{2}),
\]
\[
\bar{F}_R = (\bar{4}, 1, 2) = u^c(\bar{3}, 1, -\frac{2}{3}) + d^c(\bar{3}, 1, \frac{1}{3}) + \nu^c(1, 1, 0) + e^c(1, 1, 1) .
\]

The Higgs sector comprises the following fields. The non-trivial representations are the $SU(4)_C$ adjoint Higgs field $\Sigma = (15, 1, 1)$, the bidoublets $h, h'$ in $(1, 2, 2) + c.c.$, and the two Higgs fields

| Intersection | $SU(4)_C \times SU(2)_L \times SU(2)_R$ | $Q_C$ | $Q_{2L}$ | $Q_{2R}$ |
|--------------|---------------------------------------|-------|-----------|-----------|
| $ab$         | $3 \times F_L (4, \bar{2}, 1)$        | 1     | -1        | 0         |
| $ac$         | $3 \times \bar{F}_R (\bar{4}, 1, 2)$  | -1    | 0         | 1         |
| $ac^*$       | $H (4, 1, 2)$                         | -1    | 0         | -1        |
|              | $H (4, 1, 2)$                         | 1     | 0         | 1         |
| $aa^*$       | $S^\pm_{10} (10, 1, 1)$               | ±2    | 0         | 0         |
|              | $D^\pm_6 (6, 1, 1)$                   | ±2    | 0         | 0         |
| $cc^*$       | $\Delta_R (1, 1, 3)$                  | 0     | 0         | ±2        |
|              | $(N, \bar{N})$-singlets               | 0     | 0         | ±2        |
| $bb^*$       | $\Delta_L (1, 3, 1)$                  | 0     | ±2        | 0         |
|              | $(\nu_s, \bar{\nu_s})$-singlets      | 0     | ±2        | 0         |
| $bc$         | $h (1, 2, \bar{2})$                   | 0     | 1         | -1        |
|              | $\bar{h} (1, \bar{2}, 2)$             | 0     | -1        | 1         |
| $bc^*$       | $h' (1, 2, 2)$                        | 0     | 1         | 1         |
|              | $\bar{h}' (1, \bar{2}, 2)$            | 0     | -1        | -1        |
| $aa$         | $\Sigma (15, 1, 1)$                   | 0     | 0         | 0         |
| $bb$         | $\Sigma_{3L} (1, 3, 1)$               | 0     | 0         | 0         |
| $cc$         | $\Sigma_{3R} (1, 1, 3)$               | 0     | 0         | 0         |

Table 1: Minimal Spectrum and the corresponding quantum numbers that emerge in a D-brane configuration with $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry.
In an SO(10) embedding we note that $\bar{H}, H$ descend from the 16 and $\overline{16}$ of SO(10) respectively. In addition, a suitable set of singlet fields develop vacuum expectation values (vevs) to provide masses for the vector-like pairs and right-handed neutrinos. The minimal spectrum derived from the D-brane scenario is given in Table 1, where the transformation properties under the various quantum numbers are also shown. Notice, in particular, that the triplet pair $D_3^- + \bar{D}_3^+$ found in the decomposition of the sextets has the quantum numbers of leptoquarks since they simultaneously carry both baryon and lepton number under the definitions (3,4).

3 Two Step Spontaneous Symmetry Breaking

Having defined the spectrum of the model, in this section we proceed with the implementation of the two-step spontaneous breaking of 4-2-2 symmetry. The $SU(4)_C \rightarrow SU(3)_C \times U(1)_{B-L}$ breaking is realised with the Higgs field $\Sigma = (15, 1, 1)$ listed in Table 1:

$$\langle 15, 1, 1 \rangle \rightarrow \langle \Sigma \rangle = \begin{pmatrix} \frac{\nu_3}{\sqrt{3}} & 0 & 0 & 0 \\ 0 & \frac{\nu_3}{\sqrt{3}} & 0 & 0 \\ 0 & 0 & 0 & -\nu \end{pmatrix}.$$  \hspace{1cm} (9)

Then, at this stage, the symmetry of the model is

$$SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_L \times U(1)_R.$$  \hspace{1cm} (10)

The decomposition of the various representations of the model are shown in Table 2.

The breaking of the left-right symmetric group in (10) to the SM gauge group takes place from the non-zero vevs along the neutral components $\langle \nu_3^H \rangle, \langle \bar{\nu}_3^H \rangle$ of the right-handed doublet fields $L_H, \bar{L}_H$ in $H, \bar{H}$. Notice that these Higgs fields are also ‘charged’ under $U(1)_R$. Then, their vevs, combined with the $\bar{N}, \nu_s$ singlet non-zero vevs which will be discussed later, will also break the $U(1)_L$ and $U(1)_R$ symmetries. At the final stage, the SM symmetry breaking occurs with a non-zero vev of the bidoublet $h$.

Before studying some implications for the effective low energy theory, recall first that in the D-brane construction described above, in addition to the chiral states $F_L, \bar{F}_R$ accommodating the fermion generations, vector-like pairs such as $F^i_L, \bar{F}^i_L, F^i_R, \bar{F}^i_R, D^i_6, \bar{D}^i_6, i = 1, 2, \ldots$ etc., are usually present in the zero mode spectrum. However, such pairs arise with opposite $U(1)$ charges and, in principle, receive heavy masses of the order of the GUT scale, namely $\mathcal{W} \supset M_{GUT} (F^i_L \bar{F}^i_L + F^i_R \bar{F}^i_R + D^i_6 \bar{D}^i_6)$. For the Higgs fields given in (7,8), however, there is an additional contribution due to the coupling with the adjoint,

$$\mathcal{W}_H \supset \bar{\Pi} \Sigma H + M_H \Pi H \cdot$$  \hspace{1cm} (11)
Substituting (9), we obtain
\[ W_H \supset \left( \frac{V}{3} + M_H \right) \bar{Q}_H^c Q_H^c + (M_H - V) \bar{L}_H L_H , \] 
where \( L_H^c = (e_H^c, \nu_H^c)^T, Q_H^c = (u_H^c, d_H^c)^T \) are \( SU(2)_R \) doublets, and \( \bar{L}_H^c, \bar{Q}_H^c \) their complex conjugates. In the two step breaking pattern \( L_H^c, \bar{L}_H^c \) must develop TeV vevs to break the \( SU(2)_R \) symmetry. Therefore, choosing \( V \approx M_H \sim O(M_{\text{GUT}}) \), \( L_H^c, \bar{L}_H^c \) remain in the low energy spectrum while \( \bar{Q}_H^c Q_H^c \) acquire masses \( M_{Q_H} \sim 4\frac{3}{3}M_H \sim O(M_{\text{GUT}}) \). This “singlet-triplet” splitting is similar to the “doublet-triplet” splitting in \( SU(5) \) scenario.

4 Low Energy Phenomenology

In the previous sections we stressed that there is a natural way to provide masses for the vector-like states that might appear in the spectrum. Hence, fields such as color sextets \( D_6 \) and triplets \( \Delta_{L,R} \) of Table 1 decouple from the light spectrum as long as they appear in pairs with opposite \( U(1) \) charges. We also noticed that, when the \( SU(4)_C \) adjoint Higgs acquires its vev, the Higgs

| \( SU(3)_C \times SU(2)_L \times SU(2)_R \) | \( Q_{B-L} \) | \( Q_{2L} \) | \( Q_{2R} \) | \( B \) | \( L \) |
|---|---|---|---|---|---|
| \( 3 \times Q_L (1, 2, 1) \) | \( \frac{1}{3} \) | \( -1 \) | 0 | \( \frac{1}{3} \) | 0 |
| \( 3 \times L (1, 2, 1) \) | -1 | -1 | 0 | 0 | 1 |
| \( 3 \times Q_R^c (3, 1, 2) \) | \( -\frac{1}{3} \) | 0 | 1 | \( -\frac{1}{3} \) | 0 |
| \( 3 \times L_R^c (1, 1, 2) \) | 1 | 0 | 1 | 0 | -1 |
| \( Q_H (3, 1, 2) \) | \( \frac{1}{3} \) | 0 | 1 | \( \frac{1}{3} \) | 0 |
| \( L_H (1, 1, 2) \) | -1 | 0 | -1 | 0 | 1 |
| \( Q_H (3, 1, 2) \) | -\( \frac{1}{3} \) | 0 | -1 | \( -\frac{1}{3} \) | 0 |
| \( \bar{L}_H (1, 1, 2) \) | 1 | 0 | 1 | 0 | -1 |
| \( D_{3}^{+} (3, 1, 1) \) | \( \frac{2}{3} \) | 0 | 0 | \( \frac{2}{3} \) | 0 |
| \( \bar{D}_{3}^{+} (3, 1, 1) \) | \( \frac{2}{3} \) | 0 | 0 | \( \frac{2}{3} \) | 1 |
| \( D_{3}^{-} (3, 1, 1) \) | -\( \frac{2}{3} \) | 0 | 0 | \( -\frac{2}{3} \) | -1 |
| \( \bar{D}_{3}^{-} (3, 1, 1) \) | -\( \frac{2}{3} \) | 0 | 0 | \( -\frac{2}{3} \) | 0 |
| \( \Delta_R (1, 1, 3) \) | 0 | 0 | \( \pm 2 \) | - | - |
| \( (N, \bar{N}) \) | 0 | 0 | \( \pm 2 \) | - | - |
| \( \Delta_L (1, 3, 1) \) | 0 | \( \pm 2 \) | 0 | - | - |
| \( (\nu_s, \bar{\nu}_s) \) | 0 | \( \pm 2 \) | 0 | - | - |
| \( h, h' (1, 2, 2) \) | 0 | \( \pm 1 \) | \( \mp 1 \) | - | - |

Table 2: Spectrum of the left-right symmetric model after the breaking of the 4-2-2 symmetry.
pair $H + \bar{H}$ is a remarkable exception to this rule and, therefore, it remains in the ‘massless’ spectrum of the effective theory.

In this section we analyse the Yukawa sector and, in particular, the terms generating masses for the SM charged fields and neutrinos. The superpotential terms of the effective model should be invariant under the gauge symmetry of the original D-brane configuration of the model. We will also introduce ‘matter’ parity $\mathcal{R}$ to suppress unwanted superpotential terms and possible exotic interactions. To this end, the fields $\Sigma, H, \bar{H}, \bar{N}, h$ acquire vevs and will be assigned positive $\mathcal{R}$ parity and the fields $F_L, \bar{F}_R, N$ are assigned negative $\mathcal{R}$ parity.

Next, we analyse the contributions of the tree level terms. The third family charged fermion and Dirac neutrino mass could emerge, to a good approximation, from a common Yukawa term, $\lambda \bar{F}_R F_L h$. As a result, in the simplest constructions, the 4-2-2 symmetry imposes approximate third family Yukawa unification $\lambda_t \simeq \lambda_b \simeq \lambda_\tau \simeq \lambda_{\nu_D}$.

However, these are not the only mass terms for the neutrino sector. The following invariant couplings also involve the neutrino fields

$$\lambda \bar{F}_R h + \lambda_1 \bar{F}_R N + \lambda_2 \frac{1}{M_*} \bar{N}^2 NN,$$

where $M_*$ is the compactification scale, $M_* \gtrsim M_{\text{GUT}}$. The term $\bar{F}_R H N$ suggests that the singlet field $N = (1, 1, 1)(0, 0, -2)$ can be interpreted as a sterile neutrino. Therefore, assuming a non-zero vev for the scalar component of the $N$ singlet (positively ‘charged’ under $U(1)_R$), eq. (13) yields

$$W = \lambda_1 \langle h_u \rangle \nu^c \nu + \lambda_2 \langle H \rangle \nu^c N + \frac{\lambda_3}{M_*} \langle \bar{N} \rangle^2 NN$$

$$= m_D \nu^c \nu + V_R \nu^c N + \mu NN,$$

where we have defined

$$m_D = \lambda_1 \langle h_u \rangle, \; V_R = \lambda_2 \langle H \rangle, \; \mu = \frac{\lambda_3}{M_*} \langle \bar{N} \rangle^2.$$

The new neutral state $N$ coupled to the ordinary neutrino fields with the superpotential couplings (13) gives rise to an inverse seesaw mechanism [19, 20]. Assuming three $N$ singlet fields, in particular, we obtain a $6 \times 6$ matrix of the form

$$M_{\nu} \sim \begin{pmatrix} 0 & m_D^T & 0 \\ m_D & 0 & V_R \\ 0 & V_R & \mu \end{pmatrix}.$$

The three left handed neutrinos obtain eigenmasses of order

$$m_\nu \approx \frac{m_D^2}{V_R^2} \mu,$$

which must be sufficiently light (of order $10^{-1} - 10^{-2}$ eV) to interpret the neutrino oscillations data.
In order to provide an estimate, we make the following assumptions for the various scales involved in (17). Observing that the singlet $\bar{N}$ carries $U(1)_R$ ‘charge’, we first make the reasonable assumption that its vev is associated with the $SU(2)_R$ breaking scale
\[
\langle \bar{N} \rangle = \kappa V_R \sim O(V_R) ,
\]
where $\kappa$ is a dimensionless constant. Then, from (15) we find that the scale $\mu$ is
\[
\mu = \kappa^2 \frac{V_R^2}{M_*} .
\]
Substituting into (17), we find
\[
m_\nu \simeq \kappa^2 \frac{m_D^2}{M_*} . \tag{18}
\]
For $M_*$ of the order of the GUT scale and the Dirac mass $m_D$ of the electroweak mass scale, and with $\kappa \gtrsim 1$, we can obtain the desired left-handed neutrino masses.

The interesting fact here is that a new mass scale $\mu$ is effectively generated and it is associated with the “sterile neutrino” $N$. Since the light neutrino scale computed from the inverse see-saw matrix does not depend on the details of the $SU(2)_R$ breaking scale, we are free to choose $V_R$ (provided it satisfies the experimental bounds) so that the emergent scale $\mu$ takes the desired value. Thus, a well motivated choice is related to the possible interpretation of the dark matter puzzle through the existence of a light (order a few keV) sterile neutrino.\footnote{See for example \cite{ref1, ref2} and references therein. Also, our estimates are in agreement with recent constraints \cite{ref3} arising from from the Galactic center give a bound $m_\nu_s \leq 15$ keV.} In general, the two eigenvalues of this matrix are almost degenerate. However, with the particular choice of $SU(2)_R$ breaking scale and the existence of several neutral singlets, as is the case in the present D-brane construction, a light neutral state can emerge naturally.\footnote{See \cite{ref4} and \cite{ref5} for recent related work.}

As an example, assuming that the $\mu$ parameter in (16) represents a submatrix of the extra singlets with a scale of the order 10 keV, the scale $V_R$ is given by
\[
V_R \sim \sqrt{M_*/\text{GeV}} \cdot 10^{-3} \text{GeV} \geq 10^6 \text{GeV} .
\]
Working out the eigenvalues of the mass matrix (16) one can see that it is plausible to obtain a light neutral state of mass $O(10 \text{ keV})$ which could be interpreted as a dark matter component provided it is sufficiently long lived.

A major issue in many grand unified models is the problem of rapid proton decay. The 4-2-2 symmetry does not contain the gauge bosons that mediate dimension six baryon number violating processes. Then, the only source of proton decay is associated with dimension five or higher dimensional operators, related to graphs containing SM color triplets and other states, and the predictions are very model dependent. Operators of this kind are of the form $QQQ\ell, QQue^c, u^c d^c e^c$ etc., and may arise from non-renormalisable terms of the form $F_L F_L F_L F_L, F_L F_L F_R F_R$ and $F_R F_R F_R F_R$. Such couplings, however, are prevented from the
additional abelian symmetries $U(1)_{L,R}$ in the present construction. Then, the only available states that might mediate baryon violating diagrams, are the triplets descending from the sextets which, in principle, receive masses of order $M_{GUT}$. Non-renormalisable Yukawa couplings of the form \( \frac{1}{M_{\text{str}}} F_L F_L D_6 \nu_s \) and \( \frac{1}{M_{\text{str}}} \bar{F}_R \bar{F}_R \bar{D}_6 \tilde{N} \) give rise to diagrams similar to those discussed in [29], and their strength is determined by the vev of the scalar component of the singlet $\nu_s$. There are no constraints on the possible values of $\langle \tilde{\nu}_s \rangle$, but if this is taken to be close to the GUT scale, the proton lifetime is estimated to be of order $10^{35}-10^{36}$ yr, which will be tested by the recently approved Hyper Kamiokande experiment [30].

5 Discussion and Summary

In this letter we have analysed a class of 4-2-2 models within the framework of intersecting D-branes. The low energy spectrum contains additional particles that normally do not appear in the usual field theory based 4-2-2 models. We focused on breaking the 4-2-2 symmetry to the SM gauge group via an intermediate step $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The first stage of the two step breaking pattern is realised with the $SU(4)$ adjoint vev which does not transform under the $SU(2)_{L/R}$, and so these latter gauge symmetry factors are preserved. The second stage of symmetry breaking proceeds with the use of the $H + \bar{H} = (Q_H + L_H) + (\bar{Q}_H + \bar{L}_H)$ representations of the 4-2-2 group. Implementing a missing partner mechanism discussed in the text, the right handed doublets $\bar{L}_H^c + L_H^c$ remain in the low energy spectrum and develop TeV scale vevs along their neutral directions $\langle \nu_H^c \rangle$ and $\langle \bar{\nu}_H^c \rangle$ respectively, thereby breaking the $SU(2)_R$ symmetry.

The possibility of keV mass sterile neutrino as a potential new dark matter candidate is a striking example of this. The appearance of additional $U(1)$ symmetries which prevent rapid proton decay is certainly very helpful and a much desired feature for realistic model building. Proton decay via higher dimensional operators can yield lifetime estimates on the order of $10^{35}$-$10^{36}$ yr which will be tested by the Hyper Kamiokande experiment.

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References

[1] J. C. Pati and A. Salam, Phys. Rev. D 10 (1974) 275 Erratum: [Phys. Rev. D 11 (1975) 703].

[2] A. M. Sirunyan et al. [CMS Collaboration], Phys. Rev. D 97 (2018) no.7, 072006 [arXiv:1708.05379].

[3] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 754 (2016) 302 [arXiv:1512.01530].

[4] R. N. Mohapatra and J. C. Pati, “Left-Right Gauge Symmetry and an Isoconjugate Model of CP Violation,” Phys. Rev. D 11 (1975) 566.

[5] Q. Shafi and C. Wetterich, “Left-Right Symmetric Gauge Models and Possible Existence of a Neutral Gauge Boson with Mass in the PETRA-PEP Energy Range,” Phys. Lett. B 73 (1978) 65.

[6] R. N. Mohapatra and G. Senjanovic, “Neutrino Mass and Spontaneous Parity Violation,” Phys. Rev. Lett. 44 (1980) 912.

G. Senjanovic and V. Tello, “Origin of Neutrino Mass,” PoS PLANCK 2015 (2016) 141.

[7] G. Lazarides, Q. Shafi and C. Wetterich, “Proton Lifetime and Fermion Masses in an SO(10) Model,” Nucl. Phys. B 181, 287 (1981).

[8] J. C. Pati, Int. J. Mod. Phys. A 32, no. 09, 1741013 (2017) [arXiv:1706.09531].

[9] CMS Collaboration [CMS Collaboration], CMS-PAS-B2G-16-023.

G. Aad et al. [ATLAS Collaboration], JHEP 1512 (2015) 055 [arXiv:1506.00962].

[10] I. Antoniadis and G. K. Leontaris, “A Supersymmetric SU(4) × O(4) Model,” Phys. Lett. B 216, 333 (1989).

[11] G. K. Leontaris and J. Rizos, “A Pati-Salam model from branes,” Phys. Lett. B 510 (2001) 295 doi:10.1016/S0370-2693(01)00592-5 [hep-ph/0012255].

[12] M. Cvetic, T. Li and T. Liu, “Supersymmetric Pati-Salam models from intersecting D6-branes: A Road to the standard model,” Nucl. Phys. B 698 (2004) 163 [hep-th/0403061].

[13] M. Cvetic, D. Klevers, D. K. M. Peña, P. K. Oehlmann and J. Reuter, JHEP 1508 (2015) 087 [arXiv:1503.02068].

[14] L. E. Ibanez, F. Marchesano and R. Rabadań, JHEP 0111 (2001) 002 doi:10.1088/1126-6708/2001/11/002 [hep-th/0105155].

[15] T. P. T. Dijkstra, L. R. Huiszoon and A. N. Schellekens, Nucl. Phys. B 710 (2005) 3 [hep-th/0411129].
[16] P. Anastasopoulos, G. K. Leontaris and N. D. Vlachos, “Phenomenological Analysis of D-Brane Pati-Salam Vacua,” JHEP 1005 (2010) 011 [arXiv:1002.2937].

[17] P. Anastasopoulos, T. P. T. Dijkstra, E. Kiritsis and A. N. Schellekens, Nucl. Phys. B 759 (2006) 83 [hep-th/0605226].

[18] T. Kobayashi, S. Raby and R. J. Zhang, Nucl. Phys. B 704 (2005) 3 [hep-ph/0409098].

[19] R. N. Mohapatra and J. W. F. Valle, Phys. Rev. D 34 (1986) 1642.

[20] M. Malinsky, J. C. Romao and J. W. F. Valle, Phys. Rev. Lett. 95 (2005) 161801 [hep-ph/0506296].

[21] M. ur Rehman, V. N. Senoguz and Q. Shafi, Phys. Rev. D 75 (2007) 043522 [hep-ph/0612023].

[22] L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lust and T. R. Taylor, Phys. Rev. D 95, no. 2, 026011 (2017) [arXiv:1611.09785].

[23] S. Dodelson and L. M. Widrow, Phys. Rev. Lett. 72 (1994) 17 [hep-ph/9303287].

[24] A. Merle, V. Niro and D. Schmidt, JCAP 1403 (2014) 028 [arXiv:1306.3996].

[25] R. Yunis, C. R. Argüelles, N. E. Mavromatos, A. Moliné, A. Krut, J. A. Rueda and R. Ruffini, arXiv:1810.05756 [astro-ph.GA].

[26] K. Agashe, P. Du, M. Ekhterachian, C. S. Fong, S. Hong and L. Vecchi, Phys. Lett. B 785 (2018) 489 [arXiv:1804.06847 [hep-ph]].

[27] V. Brdar and A. Y. Smirnov, arXiv:1809.09115 [hep-ph].

[28] K. S. Babu, J. C. Pati and F. Wilczek, Nucl. Phys. B 566 (2000) 33 [hep-ph/9812538].

[29] G. K. Leontaris and Q. Shafi, Phys. Rev. D 96 (2017) no.6, 066023 [arXiv:1706.08372].

[30] K. Abe et al. [Hyper-Kamiokande Collaboration], PTEP 2018 (2018) no.6, 063C01 [arXiv:1611.06118].