Gravity induced left-right gauge theory at LHC and experimental tests

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December 2, 2014

Abstract

Embedding LHC scale $W_R, Z_R$ bosons in the non-supersymmetric GUT models of strong, weak, and electromagnetic interactions requires either a number of intermediate symmetries, or a number of nonstandard light particle degrees of freedom in the minimal chain. We show that inclusion of effect of gravity not only does away with all the additional intermediate symmetries, but also it predicts the minimal set of light Higgs scalars tailored after neutrino masses, and dilepton or trilepton signals. The heavy-light neutrino mixings are predicted and the heavy sterile Majorana fermions act as an additional source of dilepton production. Current CMS data is found to yield the bound $M_{W_R} \geq 2.5 - 2.8$ TeV. Other verifiable predictions are briefly pointed out.

1 Introduction

The standard model $SU(2)_L \times U(1)_Y \times SU(3)_C \equiv G_{213}$ partially unifies electromagnetic and weak interactions but fails to explain neutrino masses and why parity violation occurs only in weak interaction. Left-right symmetric (LRS) gauge theory $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C (g_{2L} = g_{2R}) (\equiv G_{2213D})$ predicts a number of phenomena beyond the standard model including neutrino masses and parity violation. It also goes further to suggest that the right-handed (RH) neutrino ($N$), a member of its fundamental representation, could be a heavy Majorana fermion driving seesaw mechanism for light neutrino masses and acting as a seed for baryogenesis.

As possible experimental evidence of LRS theory, it would be quite pertinent to associate these RH neutrinos to be mediating dilepton production events recently observed at the Large Hadron Collider (LHC) [4, 5] which can discriminate even if $W_R$ gauge coupling is different from the standard $W_L$ boson [6]. The minimal left-right symmetric GUT that unifies strong, weak, and electromagnetic interactions is $SO(10)$ that leaves out gravity [6]. It would be quite interesting if spontaneous symmetry breaking of non-SUSY $SO(10)$ through any one of the following two minimal symmetry breaking chains gives the LHC accessible $W_R, Z_R$ bosons

$$SO(10) \xrightarrow{M_U} G_{2213D} \text{ or } G_{2213} \xrightarrow{M_R} SM.$$  \hspace{1cm} (1)

In eq.(1) $G_{2213}$ represents the same left-right gauge theory $G_{2213D}$ but without the D-parity for which $g_{2L} \neq g_{2R}$ [6]. When the symmetry breaking proceeds through $G_{2213D}$, consistent
with extended survival hypothesis (ESH) \cite{9}, it needs large values of $\sin^2\theta_W(M_Z) \sim 0.27-0.28$ in direct conflict with electroweak precision data $(\sin^2\theta_W)_{\text{expt}} = 0.2311 \pm 0.00013$. On the other hand, solutions to renormalisation group (RG) equations consistent with this value needs an unusually larger number of nonstandard Higgs scalars and/or exotic fermions in direct violation of the ESH \cite{12}. When the symmetry breaks through $G_{2213}$, the allowed solutions also require a number of additional light particle degrees of freedom \cite{13}, although less than the $G_{2213D}$ case. In this case also the ESH has to be abandoned. Although several possibilities have been discussed earlier \cite{6,7,11}, allowed solutions for TeV scale $W_R, Z_R$ in the best identified chain of ref. \cite{7} have been noted recently to be in concordance with neutrino oscillation data \cite{10}. This model has four intermediate symmetries

$$\text{SO}(10)^{\text{MU}} \rightarrow G_{224D} \rightarrow G_{224} \rightarrow G_{2213} \rightarrow G_{2113} \rightarrow \text{SM}. \quad (2)$$

In eq.(2) $G_{224D}$ denotes the Pati-Salam symmetry $SU(2)_L \times SU(2)_R \times SU(4)_C \ (g_{2L} = g_{2R}) \ (\equiv G_{224D})$ with left-right discrete symmetry and $G_{224}$ denotes the same gauge symmetry without D-Parity. Of course a theory is finally accepted or falsified on the basis of its experimental tests and until those are carried out both the approaches to achieve LHC scale LR theory may be equally likely. \cite{6} Also in view of the LHC capability to discriminate among different models, alternative theoretical explorations for LR models with parity restoration at low scales ($g_{2L} = g_{2R}$) or high scales ($g_{2L} \neq g_{2R}$) with additional signatures would be interesting.

The purpose of this letter is to show that when effect of gravity is included through a dim.5 operator, the $SO(10)$ model gives LHC scale LR gauge theory $G_{2213}$ in the minimal symmetry breaking chain with drastically reduced size of the light Higgs spectrum that depends upon neutrino masses and the manifestation of $W_R$ through dilepton or trilepton signals. The model predicts heavy-light neutrino mixings which forms an important ingredient for multi-lepton signals at LHC especially in the $W_L$ production channel. The heavy sterile Majorana neutrinos act as additional source of dilepton signals. Other predictions of the model can be verified by charged lepton flavor violating (LFV) decay branching ratios or dominant contributions to $0\nu\beta\beta$ decay rates, or both in near future.

This letter is organized as follows. In Sec.2 we discuss the predictions of left-right gauge theory breaking scale and the grand unification scale in relation to the scale of dim.5 operator. In Sec.3 we give a short description on neutrino masses and LFV decay and in Sec.4 we discuss lepton number violation. In Sec.5 we discuss LHC signals for dilepton and trilepton production. Finally we give a brief summary of our results.

2 LHC scale LR theory

We attempt to predict the scale of LR gauge theory $G_{2213}$ in the minimal symmetry breaking chain of eq.(1).

We include the standard two-loop RG equations for gauge couplings

$$\mu \frac{\partial g_i^3}{\partial \mu} = \frac{a_i}{16\pi^2} g_i^3 + \frac{1}{(16\pi^2)^2} \sum_j b_{ij} g_j^3 g_j^2. \quad (3)$$

We also include the effect of dim.5 operator \cite{14,15}

$$\mathcal{L}_{NR} = C M_C \text{Tr}(F_{\mu\nu} \phi_{(210)} F^{\mu\nu}), \quad (4)$$

where $\phi_{210} \equiv 210_H$ Higgs representation that breaks $SO(10) \rightarrow G_{2213}$ at the GUT scale by acquiring vacuum expectation value (VEV) along its $G_{2213}$ singlet direction and $M_C$ is the scale

\footnote{Compatibility of such LR theories \cite{12,13} with neutrino oscillation data is yet to be investigated.}
of the dim.5 operator\[3\] For the gauge kinetic field tensor we have

\[ F_{\mu\nu} = \partial_{\mu}W_{\nu} - \partial_{\nu}W_{\mu} + ig[W_{\mu}, W_{\nu}], \]

\[ W_{\mu} = \frac{1}{4} \sum_{i,j=1}^{10} \sigma^{ij}W^{ij}_{\mu}, C = -\frac{3}{8} \]

where \( \frac{1}{2}\sigma^{ij}(W^{ij}_{\mu}) \), and i,j=1,2,3......10 denote the 45 generators (gauge bosons) of SO(10). The GUT-scale boundary conditions are modified by the dim.5 operator

\[ \alpha_{2L}(M_U)(1 + \epsilon_{2L}) = \alpha_{2R}(M_U)(1 + \epsilon_{2R}) = \alpha_{BL}(M_U)(1 + \epsilon_{BL}) = \alpha_{3C}(M_U)(1 + \epsilon_{3C}) = \alpha_G, \]

where the \( \epsilon_{i} \) terms arise due to the dim.5 operator and \( \alpha_G \) is the effective GUT fine structure constant. The resulting analytic formulas for the unification mass \( M_U \) and the LR scale \( M_R \) are \[16\]

\[ \log \frac{M_R}{M_Z} = \frac{1}{(XZ' - X'Z)}[(XP_s - X'P_b) + (X'P_2 - X'\Sigma_2) - \frac{2\pi}{\alpha_G}(X\epsilon'' - X'\epsilon')], \]

\[ \log \frac{M_U}{M_Z} = \frac{1}{(XZ' - X'Z)}[(ZP_b - ZP_s) + (Z\Sigma_2 - Z'\rho_2) - \frac{2\pi}{\alpha_G}(Z'\epsilon' - Ze'')], \]

The terms on the right-hand side reduce to the usual renormalisable ones in the limit \( \epsilon' = \epsilon'' = 0 \).

Various notations occurring in eq.(7) and eq.(8) are

\[ P_s = \frac{2\pi}{\alpha(M_Z)} \left( 1 - \frac{8}{3} \frac{\alpha(M_z)}{\alpha_s(M_z)} \right), \]

\[ P_b = \frac{2\pi}{\alpha(M_Z)} \left( 1 - \frac{8}{3} \sin^2 \theta_W(M_z) \right), \]

\[ X = a'_{2R} + \frac{2}{3}a'_{BL} - \frac{5}{3}a'_{2L}, \]

\[ Z = \frac{5}{3}(a_Y - a_{2L}) - \left( a'_{2R} + \frac{2}{3}a'_{BL} - \frac{5}{3}a'_{2L} \right), \]

\[ X' = a'_{2R} + \frac{2}{3}a'_{BL} + a'_{2L} - \frac{8}{3}a'_{3C}, \]

\[ Z' = \frac{5}{3}a_Y + a_{2L} - \frac{8}{3}a_{3C} - \left( a'_{2R} + \frac{2}{3}a'_{BL} + a'_{2L} - \frac{8}{3}a'_{3C} \right), \]

\[ \rho_2 = 2\pi \left[ \theta'_{2R} + \frac{2}{3}\theta'_{BL} - \frac{5}{3}\theta'_{2L} + \frac{5}{3}(\theta_Y - \theta_{2L}) \right], \]

\[ \Sigma_2 = 2\pi \left[ \theta'_{2R} + \frac{2}{3}\theta'_{BL} + \theta'_{2L} - \frac{8}{3}\theta_{3C} + \frac{5}{3}\theta_Y + \theta_{2L} - \frac{8}{3}\theta_{3C} \right], \]

\[ \epsilon' = \epsilon_{2R} + \frac{2}{3}\epsilon_{BL} - \frac{5}{3}\epsilon_{2L}, \]

\[ \epsilon'' = \epsilon_{2L} + \epsilon_{2R} + \frac{2}{3}\epsilon_{BL} - \frac{8}{3}\epsilon_{3C}. \]

The first, second, and third terms in the R.H.S. of eq.(7) and eq.(8) represent one loop, two-loop and gravitational effects respectively. In particular the combined GUT symmetry breaking VEVs of \(< \eta(1,1,1) >, < \eta'(1,1,15) > < 210_H > \) can be written as \[15\]

\[ ^3\text{When} M_G \simeq M_{Planck}, \text{we call it a gravitational effect in four-dimensional space time but when} M_G < M_{Planck} \text{such a term may arise due to compactification of extra dimensions of Kaluza-Klein type} \[14\] \]
\[
\langle \phi_{(210)} \rangle = \frac{\phi_0}{8\sqrt{2}} (-\Gamma_1 \Gamma_2 \Gamma_3 \Gamma_4 + \Gamma_1 \Gamma_2 \Gamma_5 \Gamma_6 + \Gamma_3 \Gamma_4 \Gamma_5 \Gamma_6) + \Gamma_7 \Gamma_8 \Gamma_9 \Gamma_{10},
\]

(10)

where we have used \( \Phi_{(210)} = \frac{1}{2} \Gamma_j \Gamma_k \Gamma_l \times \phi^{jkl} \), \( < \Phi_{1234} > = < \Phi_{1256} > = < \Phi_{3456} > = < \Phi_{78910} > \) leading to

\[
\epsilon_{2R} = -\epsilon_{2L} = -\epsilon_{3C} = \frac{1}{2} \epsilon_{BL} = \epsilon
\]

\[
\epsilon = -\frac{C M_D}{2M_C} \left( \frac{3}{2\pi G} \right)^2.
\]

(11)

**Heavy pseudo Dirac neutrinos**

In this case 210_H breaks SO(10) and \( D \) – Parity to \( G_{2213} \) which further breaks to SM by the doublet \( \chi_R(1,2,-1,1) \subset 16_H \). The SM theory breaks to the low-energy symmetry by the standard Higgs doublet \( h(2,1,1) \subset \phi(2,2,0,1) \subset 10_H \). With such minimal Higgs content, respective beta function coefficients are presented in Table 1. Three additional singlet fermions \( (S_i, i = 1, 2, 3) \), one for each generation are added in case of \( SO(10) \) theory\(^4\). The VEV of the RH doublet \( \chi_R \) also generates the \( N - S \) mixing mass term \( M = Y_{\chi} < \chi_R > \) which, along with the Dirac neutrino mass \( M_D \), is known to contribute to leptonic non-unitarity, leptonic CP violation, and lepton flavor violation. After block diagonalisation of the 9 \( \times \) 9 neutral fermion mass matrix

\[
\mathcal{M} = \begin{pmatrix}
0 & M_D & 0 \\
M_D^T & 0 & M \\
0 & M^T & \mu_S
\end{pmatrix}
\]

(12)

the RH neutrinos form three pairs of heavy pseudo-Dirac neutrinos with \( |M_{N_i}| \simeq |M_i| (i = 1 - 3) \). In combination with \( M_D, M \) and the global lepton number violating fermion singlet mass term \( \mu_S \), the model drives the inverse seesaw mechanism for neutrino masses \( [18, 23] \)

\[
m_\nu = \frac{M_D}{M} \mu_S \left( \frac{M_D}{M} \right)^T,
\]

(13)

where, for simplicity, we use \( M = \text{diag}(M_1, M_2, M_3) \).

Using the numerical values of the available data \( \sin^2 \theta_W(M_Z) = 0.23116 \pm 0.00013 \), \( \alpha(M_Z) = 1/127.9 \) and \( \alpha_{\text{QCD}}(M_Z) = 0.1184 \pm 0.0007 \) in eq.\(^7\), eq.\(^8\), we obtain solutions to the values of \( M_R, M_U, \sigma \) and \( M_C \) which are presented in Table 2. It is clear that the gravity induced solutions as low as \( V_{\chi n} \sim M_R \simeq 9 - 25 \text{ TeV} \) are allowed and the model predicts the \( W_R \) mass \( M_{W_R} \simeq g_2 R V_{\chi n} \sim 3 - 10 \text{ TeV} \) in the case of minimal combination of the light Higgs sector with only two doublets, \( D_\phi = D_\chi = 1 \). The scale of dim.5 operator turns out to be almost at the reduced Planck scale signifying that the induced nonrenormalisable corrections could be genuinely due to gravitational effects of four dimensional space time. The estimated proton lifetime for the \( p \rightarrow e^+ \pi^0 \) mode is found to be beyond the range accessible to Super-Kamiokande or Hyper-Kamiokande collaborations. It is interesting to note that \( \sigma = \frac{\sigma_{\text{LS}}}{\sigma_{\text{27}}} \simeq 1.27 \) is only 27% larger compared to the value \( \sigma = 1 \) in the manifest LRS model \(^2\).

**Heavy RH Majorana neutrinos**

For this purpose, in addition to the Higgs representations of Model-I, we require the \( SO(10) \) representation \( 126_H \supset \Delta_R(1,3,-2,1) \) under \( G_{2213} \) that carries \( B - L = -2 \) with corresponding coefficients given in Table 1. When the RH triplet acquires VEV \( \langle \Delta_R^0 \rangle = V_{\Delta_R} \), \( G_{2213} \) symmetry

\(^4\text{In } E_6 \text{ theory each singlet fermion is part of its fundamental representation which decomposes under } SO(10) \text{ as } 27 = 16 + 10 + 1 \).
Table 2: Predictions for LR scale $M_R$ and GUT scale $M_U$ in relation to $M_C$, the scale of dim.5 operator.

| Model      | $\epsilon$ | $a_G^{-1}$ | $M_R$(GeV) | $M_U$(GeV) | $C = \frac{\alpha_C}{\sqrt{2}}$ | $\sigma$ | $M_C$(GeV) | $C = -\frac{\epsilon}{2}$ |
|------------|-------------|------------|------------|------------|----------------------------------|----------|------------|----------------------------------|
| Model-I    | 0.082       | 49.6       | $2.32 \times 10^4$ | $2.77 \times 10^{16}$ | $1.274$ | $1.12 \times 10^{18}$ | $-1.36$ |
|            | 0.083       | 49.6       | $1.9 \times 10^4$ | $2.75 \times 10^{16}$ | $1.277$ | $1.1 \times 10^{18}$ | $-1.37$ |
| Model-II   | 0.044       | 49.6       | $6.66 \times 10^4$ | $2.89 \times 10^{16}$ | $1.20$ | $2.37 \times 10^{18}$ | $-1.35$ |
|            | 0.042       | 49.7       | $4.22 \times 10^4$ | $2.84 \times 10^{16}$ | $1.212$ | $2.2 \times 10^{18}$ | $-1.33$ |
|            | 0.045       | 49.5       | $2.28 \times 10^4$ | $2.77 \times 10^{16}$ | $1.238$ | $1.98 \times 10^{18}$ | $-1.32$ |

is broken down to SM and RH neutrinos acquire heavy masses through Yukawa interaction term $\mathcal{H}_{16,16,126}_H$ leading to $M_N = fV\Delta_n$ that replaces the central part of the $3 \times 3$ null matrix of eq. (12). The RH doublet $\chi_R(1, 2, -\frac{1}{2}, 1)$ apart from taking part in symmetry breaking process rather weakly, generates the N-S mixing mass term $M$ as noted in the case of Model-I leading to gauged inverse seesaw formula for neutrino masses provided $M_N > M > M_D, \mu_S$, a condition well known in extended seesaw mechanism [22]. The would-be dominant type-I seesaw term in this model cancels out in such decoupling limit [10,17,21] leading to gauged inverse seesaw formula of eq.(13) to explain the neutrino oscillation data. There are two heavy Majorana neutrino mass matrices: $m_N$ for RH neutrino and $m_s$ for sterile neutrino under the constraint $m_N >> m_s$

$$m_s = \mu_s - M \frac{1}{M_N} M^T + ...$$  \hspace{1cm} (14)

The heavy RH Majorana neutrino mass matrix is very close to its gauged value,

$$m_N = M_N + ...$$  \hspace{1cm} (15)

These two types of heavy Majorana neutrinos emerging as outcome of the extended seesaw mechanism mediate neutrino-less double beta decay in the $W_L - W_L, W_L - W_R,$ and $W_R - W_R$ channels. Further both of them are capable of mediating the dilepton production process.

Using numerical values of $a_i$ and $b_{ij}$ in eq.(7), eq.(8), and eq.(9) and following the same procedure as outlined for Model-I, the solutions for mass scales $M_R, M_U, M_C,$ and $\sigma$ are also presented for this Model-II in Table 2. It is clear that in this case low-mass RH gauge bosons are also permitted at the LHC energy scale for $\sigma \simeq 1.2$ and $M_C$ values almost at the reduced Planck scale which leads to the interpretation that the additional corrections could be due to gravity effects.
3 Neutrino masses and lepton flavor violation

The Dirac neutrino mass matrix $M_D$ occurring in eq. (13) is determined by the GUT-scale fitting of the extrapolated values of all charged fermion masses obtained by following the bottom-up approach [19] and running it down to the TeV scale following top-down approach as explained in the corresponding cases [10,23,24]. While the procedure followed in ref. [23] is used for Model-I, the procedures followed in ref. [10,24] is utilized for Model-II. An additional bidoublet $\phi' \subset 10_H'$ is needed to fit fermion masses without affecting coupling unification substantially. The Higgs bidoublet $\xi (2,2,15) \subset 126_H$ acquires the induced VEV $v_\xi \simeq 10 - 50$ MeV [20] which, along with the direct VEVs of the two bidoublets, enables fitting all charged fermion masses in Model-II. A byproduct of this fitting is the diagonalised version of the heavy RH neutrino mass matrix $\hat{M}_N = \hat{f}V_R = \text{diag}(\hat{M}_{N_1}, \hat{M}_{N_2}, \hat{M}_{N_3})$. (16)

where, in our Model-II,

$\hat{M}_{N_1} \simeq 150 \text{GeV} - 1.5 \text{TeV}$,

$\hat{M}_{N_2} \simeq 500 \text{GeV} - 2.0 \text{TeV}$,

$\hat{M}_{N_3} \simeq 2.0 \text{GeV} - 7.5 \text{TeV}$. (17)

In the absence of $126_H$ in Model-I, the dim.6 operator $F_{ij}^{16_i, 16_j, 10_H, 45_H}/M^{12}$ discharges the equivalent role where $M' \sim M_{Planck}$. This $45_H$ remains near the GUT scale without affecting the particle spectrum at the LHC scale. Up to a good approximation, in both the models the Dirac neutrino mass matrix at the LHC scale is

$$M_D(M_R^0) = \begin{pmatrix}
0.0151 & 0.0674 - 0.0113i & 0.1030 - 0.2718i \\
0.0674 + 0.0113i & 0.4758 & 3.4410 + 0.0002i \\
0.1030 + 0.2718i & 3.4410 - 0.0002i & 83.450
\end{pmatrix} \text{GeV.}$$ (18)

The dominant source of LFV is through the $W_L$-loop in both the models and there are two types of heavy Majorana Fermion exchange contributions in case of Model-II. The RH neutrino exchange contribution can be considered subdominant since $M_{N_i} \gg M_i$. Using the relevant analytic formulas [21] we estimate LFV decay branching ratios $\mu \rightarrow e + \gamma, \tau \rightarrow e + \gamma$ and $\tau \rightarrow \mu + \gamma$ as shown in Table 3 where the allowed values of $M_i,(i = 1, 2, 3)$ satisfying the non-unitarity constraints have been also given [10,24]. As the predicted values are $3 - 5$ orders smaller than the current experimental limits they may be accessible to ongoing or planned searches with improved accuracy.

| $M (\text{GeV})$ | $BR(\mu \rightarrow e\gamma)$ | $BR(\tau \rightarrow e\gamma)$ | $BR(\tau \rightarrow \mu\gamma)$ |
|-----------------|-------------------------------|-------------------------------|-------------------------------|
| $(50, 200, 1711.8)$ | $1.19 \times 10^{-16}$ | $4.13 \times 10^{-15}$ | $5.45 \times 10^{-15}$ |
| $(100, 100, 1286)$ | $1.07 \times 10^{-15}$ | $2.22 \times 10^{-14}$ | $2.64 \times 10^{-12}$ |
| $(100, 200, 1702.6)$ | $1.14 \times 10^{-16}$ | $4.13 \times 10^{-15}$ | $5.52 \times 10^{-13}$ |
| $(500, 500, 1140)$ | $1.35 \times 10^{-16}$ | $1.25 \times 10^{-14}$ | $1.77 \times 10^{-12}$ |

Using the set of values on $M$ from Table 3 and the Dirac neutrino mass matrix from eq. (13), we fit the available data on neutrino masses and mixings through inverse seesaw formula of eq. (13) for all the three types of mass hierarchies: NH, IH, and QD. In each case the fit gives a set of

\[ Any additional SO(10) Higgs representations or higher dimensional operators which may be needed for charged fermion mass fits at the GUT scale do not affect the LHC scale particle spectrum. \]
elements for $\mu_S$. As an example by using $M = \text{diag.(500, 500, 1140)}$ GeV and the QD light neutrino hierarchy of common mass $m_0 = 0.2$ eV \[23\] gives

$$
\mu_S(\text{GeV}) = \begin{pmatrix}
0.9932 - 0.0124i & -0.1908 + 0.0022i & 0.0066 - 0.0033i \\
-0.1908 + 0.0022i & 0.0370 - 0.0004i & -0.0013 + 0.0006i \\
0.0066 - 0.0033i & -0.0013 + 0.0006i & 0.0003 - 0.00004i
\end{pmatrix} \text{GeV} \tag{19}
$$

Different aspects of LFV in non-SUSY $SO(10)$ have been discussed in ref. \[10, 23, 24\] and our predictions in the corresponding cases are similar.

## 4 Lepton number violation

The standard contribution in the $W_L - W_L$ channel is due to light neutrino exchanges. But, one important aspect of this Model-II is that even in the $W_L - W_L$ channel, the singlet fermion exchange allowed within the extended seesaw mechanism, can yield much more dominant contribution to $0\nu\beta\beta$ decay rate than the standard one. The contributions due to the exchanges of heavy $W_R, \Delta^+_R$ and the RH neutrino in the extended seesaw framework are negligible compared to the light neutrino and singlet fermion exchange contributions in the $W_L - W_L$ channel. The leading order contributions have been summarized in in the Table 4.

Table 4: Formulas for amplitudes and effective mass parameters in $W_L - W_L$ channel for $0\nu\beta\beta$ decay.

| Channel | Amplitude | Effective mass parameter |
|---------|-----------|-------------------------|
| $W_L - W_L$ | $\nu$ | $A^{\nu L} \propto \frac{g^4}{M_{W_L}^2} \sum_{i=1,2,3} \frac{(V_{\nu i})^2 m_{\nu i}}{m^2_{\nu i}}$ |
| | $S$ | $A^S \propto \frac{g^4}{M_{W_L}^2} \sum_{j=1,2,3} \frac{(V_{\nu j})^2 m_{S j}}{m^2_{S j}}$ |
| | $N$ | $A^N \propto \frac{g^4}{M_{W_L}^2} \sum_{k=1,2,3} \frac{(V_{\nu k})^2 m_{N k}}{m^2_{N k}}$ |

Using expressions in Table 4 in the combined contribution of light and sterile neutrinos to the effective mass parameter

$$m_{\nu eff} = m^{ee.L}_\nu + m^{ee.S}_S,$$ \tag{20}

the first term on the RHS is a function of light neutrino masses, PMNS mixings, and phases whereas, in the second term, $m_{S_i}, i = 1,2,3$ are mass eigen values of the sterile neutrinos and
\[ \nu_{ei}^S = \frac{(M_D/M)_{ei}}{(M_D)_{ei}/M_i}. \]  
Whereas the first term is known from neutrino oscillation data, the second term is known from Table (4), eq. (18), the \( M_i \) values of Table 3, and the computed values of \( m_{S_i} \) through eq. (14) and eq. (17).

Out of various contributions in the \( W_L-W_L \) channel analysed in detail [24], we give an example of scattered plot of predicted half life as a function of the lightest sterile neutrino mass eigen value in Fig. 4 in the normally hierarchical (NH) case of the light neutrino masses.

Saturation of current experimental bound on \( 0\nu\beta\beta \) decay half life gives the lower bound on the lightest sterile neutrino mass:

\[ m_{S_1} \geq 17 \pm 3 \text{ GeV}. \]  
In a similar manner bounds evaluated in the IH and QD cases have been found to be consistent with this value.

## 5 LHC signals of heavy pseudo Dirac and Majorana neutrinos

At the LHC, the parton-level generation of a heavy neutrino can be realized in the following way

\[ q\bar{q} \rightarrow W_L/W_R \rightarrow l^+N(l^-N), (l = e, \mu, \tau) \]  
provided this process is kinematically feasible. This has lepton-number conserving (LNC) or lepton number violating (LNV) decay modes depending on whether \( N \) is pseudo Dirac as in Model-I or Majorana as in Model-II. We use the parton level differential cross section [28]

\[
\frac{d\sigma_{LHC}}{d\cos \theta} = \frac{k\rho}{32\pi s} \frac{\hat{s} + M^2}{\hat{s}} \frac{\theta}{48} \left( \hat{s} - m_W^2 \right)^2 + m_W^4 \Gamma_W^2 (2 + \rho \cos^2 \theta) \]  
where \( k = 3.89 \times 10^8 \text{pb}, \hat{s} \) is the square of centre-of-mass energy of the colliding partons, \( M \) is mass of \( N \), and \( \rho = (\hat{s} - M^2)/\hat{s} + M^2 \). The total production cross section at the LHC is

\[
\sigma_{prod} = \int d\sqrt{s} \int d\cos \theta \int d\tau \frac{A_\tau}{x} \left[ f_u(x, Q) f_d(\tau, Q) \frac{d\sigma_{LHC}}{d\cos \theta} + (u \rightarrow \bar{u}, d \rightarrow d) \right], \]  
where \( \tau = \hat{s}/E_{CM}^2 \) and \( E_{CM} \) is centre-of-mass energy of the LHC. The Feynman diagrams for trilepton(dilepton) production mechanism is shown in the left-panel (right-panel) of Fig.2

### Trilepton signals

The RH neutrino in Model-I being pseudo Dirac can not mediate like-sign dilepton production and the best channel for the trilepton mode is where \( W_L/W_R \) decays to leptonic final states:

\[ pp \rightarrow W_R^\pm \rightarrow N l^\pm \rightarrow W_L/W_R^{\pm}\ell^\pm \rightarrow \nu l^\pm\ell^\pm \ell^\pm \]  

Figure 2: Feynman diagrams for trilepton (left-panel) and dilepton (right-panel) signals at the LHC in various channels: LL, RR, LR, and RL where instead of confining to the exchange of real \( W_R \) boson [20], the general possibilities including both real and virtual \( W_L, W_R \) exchanges [27] in the second stage have been included.

The inclusive cross-section for the trilepton state in a generic seesaw model is given by [28]

\[
\sigma(pp \rightarrow l_1 l_2 l_3 + T_{me}) = \sigma_{prod}(pp \rightarrow W^* \rightarrow N l_1) Br(N \rightarrow l_2 W) Br(W \rightarrow l_3 \nu). \]  

8
Here, \(T_{mc}\) stands for the missing transverse energy and the \(W_L\) branching ratio \(\text{Br}(W \to l \nu) = 0.21\) \cite{29}. We have assumed \(m_N > m_W\). Although this condition is needed for kinematic feasibility of the decay \(N \to l_2 W \to l_2 l_3 \nu\) when the \(W\) is the real intermediate boson of the SM, this is not required for virtual \(W^*\) exchange to give \(N \to l_3 W^* \to l_2 l_3 \nu\). One important aspect of this model is that the fermion mass fitting and LFV constraint predicts all the elements of the heavy-light neutrino mixing matrix \(V_{lS} = \frac{\Sigma}{\mu}\). For example using eq. \(18\) and \(M_{N_2} \simeq M_2 = 50\) GeV, the light-heavy neutrino mixing parameter is \(|V_{lN_{N_1}}|^2 = 9.8 \times 10^{-5}\). Thus the heavy-light neutrino mixing is determined and varies inversely as the corresponding heavy pseudo Dirac neutrino mass exchanged.

For computation of the production cross section we have utilized the MRST parton distribution functions \cite{31} in eq. \(23\). Using our ansatz for heavy light neutrino mixing matrix \(M_D/M\) in the pseudo Dirac case, eq. \(22\), eq. \(23\), and eq. \(24\), our predicted results on trilepton signals in the \(LL\) channel are shown for LHC energy \(\sqrt{s} = 14\) TeV in Fig.3 where \(l_3, l_3^* = e^\pm e^\mp\) and \(l_3 = e^\pm\) or \(\mu^\pm\) in the upper blue curve and the mediating heavy fermion is the pseudo Dirac \(N_2\). The corresponding trilepton signal as a function of the pseudo Dirac mass \(M_{N_1}\) is shown as lower red curve in the same figure for which \(l_3^* = \mu^\pm \mu^\pm\) but \(l_3 = e^\pm\) or \(\mu^\pm\).

![Figure 3: Signal cross sections for trilepton final states for the in the LL channel as a function of heavy pseudo Dirac mass \(M_{N_1}\) (red curve) and \(M_{N_2}\) (blue curve) at \(\sqrt{s} = 14\) TeV.](image)

At \(\sqrt{s} = 14\) TeV, the predicted trilepton signal cross sections in the \(WR - WR\) channel are shown in Fig.\(4\)(a) when \(l_3^*, l_3 = e^\pm e^\mp\) and \(l_3 = e^\pm\) or \(\mu^\pm\). In Fig. \(4\)(b) the predicted signal cross sections are for \(l_3 = \mu^\pm \mu^\pm\), and \(l_3 = e^\pm\) or \(\mu^\pm\) also in the same channel. In this case

![Figure 4: Signal cross sections for trilepton final states for the RR channel at \(\sqrt{s} = 14\) TeV.](image)

**Dilepton signals**

The Model-II has two types of heavy Majorana neutrinos. In this case, the collider signatures of heavy RH neutrino and the sterile neutrino can manifest at LHC energies through dilepton production in various channels of which we have examined the \(WL - WL\) and \(WR - WR\), channels. Thus compared to other models \cite{26,27} this model gives an additional source of dilepton production especially in \(LL, LR\) and \(RL\) channels due to heavy sterile neutrino exchanges.

The signal cross-section for the production of the RH neutrino or sterile neutrino including the real or virtual \(WL, WR\) exchanged at the second stage is given by

\[
\sigma(pp \to Nl^\pm \to l^\pm l^\pm jj) = \sigma_{\text{prod}}(pp \to WL, WR \to Nl^\pm) \times \text{Br}(N \to l^\pm jj) \tag{25}
\]
where the branching ratio

$$Br(N \rightarrow l^\pm jj) = \frac{\Gamma(N \rightarrow l^\pm jj)}{\Gamma_{tot}^N} \times Br(W \rightarrow jj)$$

(26)

here $Br(W \rightarrow jj) = 0.676$ [29]. For heavy Majorana neutrino exchange, our results are shown for $\sqrt{s} = 14$ TeV with MRST parton distribution functions in Fig. 5 and Fig. 6(a) respectively, in the LL and RR channels. For sterile neutrino exchanges, our computed results are shown in Fig. 7 and Fig. 8, respectively, in the LL and RR channels.

Figure 5: Signal cross sections for dilepton final states in the LL channel at $\sqrt{s} = 14$ TeV. The middle curve in the right panel represents the predicted signal cross section of our Model-II while the curves above and below the middle one represent signal cross sections of two benchmark scenarios as [27].

Figure 6: Left-panel: Predictions for dimuon signal cross section at $\sqrt{s} = 14$ TeV in the RR channel as a function of $M_{N_2}$. Right-panel: Comparison of different models, (a) manifest LRS model (green small dashed), (b) this analysis of Model-II (solid red), and (c) ref. [10] (magenta long-dashed).

At 30 fb$^{-1}$ luminosity, the number of events for heavy neutrino mass $M_N=100$ GeV are 22.17 for $W_L - W_L$ channel and 1380 for $W_R - W_R$ channel respectively. Here, it is noteworthy that the dilepton signal is more dominant over trilepton signal in the $W_R - W_R$ channel. It is also found that our estimated signal cross section for heavy neutrino in $W_L - W_L$ channel lies in between those of two benchmark scenarios as discussed in ref. [27]. Hence this model predicts a signal cross section having a better probability of observing heavy neutrinos at the LHC or the other ongoing experiments.

Using $\sqrt{s} = 8$ TeV we have shown our model predictions of the Di electron and dimuon signal cross sections in the left-panel and the right panel, respectively in Fig. 9 for $W_R$ production, and compared them with the CMS data [5]. The line I with uncertainty band is the prediction of the manifest LR model ($\sigma = 1$). The line II is our Model-II prediction for $(g_{2L}/g_{2R})^2 = \sigma = 1.3$ and
Figure 7: Signal cross sections mediated by sterile Majorana neutrinos for dilepton final states in the $LL$ channel at $\sqrt{s} = 14$ TeV.

$V_{e1}^2 = 1$ in both the panels whereas the line III represents our prediction in the same model for the same value of $\sigma$ but $V_{e1}^2 = 0.8$ (left-panel) and $V_{e2}^2 = 0.7$ (right-panel). These mixings among the RH neutrinos are similar to the PMNS mixings for light neutrinos. The possibility of trilepton signals by heavy Majorana fermion exchanges [30] would be presented elsewhere.

Figure 8: Same as Fig. 7 but in the RR channel.

Figure 9: Predictions of dielectron (left-panel) and dimuon (right-panel) signal cross section (lines II and III) for $W_R$ production at $\sqrt{s} = 8$ TeV and their comparison with the LHC data for which the green (red) band is the 1$\sigma$ (2$\sigma$) limit. The zig-zag dotted (solid) curve represents expected (observed) results of measurements. The line III with spreaded uncertainty is the prediction of manifest LRS model $g_{2L} = g_{2R}$ [3].
Summary: In summary we have shown the realization of LHC scale LR gauge theory in the minimal chain with minimal light Higgs spectrum consistent with neutrino oscillation data and all charged fermion mass fit in non-SUSY SO(10) model. The content of the light Higgs spectrum depends upon the pseudo-Dirac (Majorana) nature of the heavy RH neutrino that predicts trilepton (like-sign dilepton) signal for $W_R$ production at the LHC. Only for gauge coupling unification in the pseudo Dirac case the Model-I has just one bidoublet and one RH doublet carrying $B-L=-1$ and one more bidoublet-doublet when all fermion mass fit is desired. In the heavy RH Majorana neutrino case, in addition to the light Higgs of Model-I in the corresponding cases, the Model-II needs just one more RH triplet Higgs scalar carrying $B-L=-2$. We have reported predictions of trilepton (dilepton) signal cross sections in the Model-I (Model-II) in the $LL$ and $RR$ channels at $\sqrt{s} = 14$ TeV and also predicted new signal cross sections due to the mediation of heavy sterile neutrinos in Model-II. The heavy-light neutrino mixings are predicted because of the determined value of the Dirac neutrino mass matrix and LFV constraints. The model predicts $g_{2L}^2/g_{2R}^2 = 1.2 - 1.3$ leading to the lower bound $M_{W_R} \geq 2.5 - 2.8$ TeV from the CMS data. In addition, the model also predicts LFV branching ratios and dominant $0\nu\beta\beta$ decay rate accessible to ongoing experiments even when light neutrino masses are normally hierarchical.

Acknowledgment: M. K. P. thanks Department of Science and Technology, Govt. of India for the research project SB/52/HEP-011/2013. B. S. thanks SOA University for a research fellowship. The authors thank Ram Lal Awasthi and Samiran Bose for computational help.

References

[1] J. C. Pati and A. Salam, Phys. Rev. D 8 (1973) 1240; ibid. D 10 (1974) 275.

[2] R. N. Mohapatra and J. C. Pati, Phys. Rev. D 11 (1975) 2558; Phys. Rev. D 11 (1975) 566; G. Senjanovic and R. N. Mohapatra, Phys. Rev. D 12 (1975) 1502.

[3] R. N. Mohapatra and G. Senjanovic, Phys. Rev. D 23 (1981) 165.

[4] CMS Collaboration, S. Chatrchyan et al., Phys. Rev. Lett. 109 (2012).

[5] CMS Collaboration, V. Khachatryan et al., arXiv:1407.3683 [hep-ex]

[6] D. Chang, R. N. Mohapatra, and M. K. Parida, Phys. Rev. Lett. 52 (1984) 1072; D. Chang, R. N. Mohapatra, and M. K. Parida, Phys. Rev. D 30 (1984) 1052.

[7] D. Chang, R. N. Mohapatra, J. Gipson, R. E. Marshak, and M. K. Parida, Phys. Rev. D 31 (1985) 1718.

[8] P. S. Bhupal Dev and R. N. Mohapatra, Phys. Rev. D 81 (2010) 013001.

[9] F. del Aguila, L. Ibanez, Nucl. Phys. B 177 (1981) 60; R. N. Mohapatra, G. Senjanovic, Phys. Rev. D 27 (1983) 1601.

[10] R.L. Awasthi, M.K. Parida, and S. Patra, JHEP 1308 (2013) 122.

[11] S. Bertolini, T. Schwetz, M. Malinsky, Phys. Rev. D 73 (2006) 115012; S. Bertolini, Luca Di Luzio, M. Malinsky, Phys. Rev. D 80 (2009) 015013.

[12] M. Lindner, M. Weiser, Phys. Lett. B 383 (1996) 405.

[13] C. Albaez, M. Hirsch, M. Malinsky, J. C. Romao, Phys. Rev. D 89 (2014) 035002.
[14] Q. Shafi, C. Wetterich, Phys. Rev. Lett. (1984); C. T. Hill, Phys. Lett. B (1984); P. Langacker, N. Polonsky, Phys. Rev. D 47 (1993) 4028; L. J. Hall and U. Sarid, Phys. Rev. Lett. 70 (1993) 26; M. K. Parida, P. K. Patra and A. K. Mohanty, Phys. Rev. D 39 (1989) 316; S. K. Majee, M. K. Parida, and A. Raychaudhuri, Phys. Lett. 668 (2008) 299.

[15] M.K. Parida, P.K. Patra, Phys.Lett. B 234 (1990) 45.

[16] M. K. Parida, B. Purkayastha, C. R. Das, B. D. Cajee, Eur. Phys. J. C 28 (2003) 353.

[17] S. K. Kang and C. S. Kim, Phys. Lett. B 646 (2007) 248; J. Ellis, D. V. Nanopoulos and K.Olive, Phys. Lett. B 300 (1993) 121; S. K. Majee, M. K. Parida, and A. Raychaudhuri, ref. [14].

[18] R. N. Mohapatra, Phys. Rev. Lett. 56 (1986) 61; R. N. Mohapatra and J. W. F. Valle, Phys. Rev D 34 (1986) 1642.

[19] C. R. Das and M. K. Parida, Eur. Phys. J. C 20 (2001) 121.

[20] K. S. Babu and R. N. Mohapatra, Phys. Rev. Lett. 70 (1993) 2845.

[21] A. Ilakovac and A. Pilaftsis, Nucl. Phys. B 437 (1995) 491.

[22] W. Grimus and L. Lavoura, JHEP, 0011 (2000) 042; M. Mitra, F. Vissani, G. Senjanovic, Nucl. Phys. B 856 (2012) 26.

[23] R.L. Awasthi, M.K. Parida, Phys. Rev. D 86 (2012) 093004.

[24] M. K. Parida, R. L. Awasthi, P. K. Sahu, [arXiv:1401.1412[hep-ph]].

[25] J. Barry, L. Dorame and W. Rodejohann; Eur. Phys. J. C 72 (2012) 2023; J. Barry and W. Rodejohann , JHEP, 1309 (2013) 153.

[26] W.Y. Keung and G. Senjanovic, Phys. Rev. Lett. 50 (1983) 1427.

[27] Chien-Yi Chen, P. S. Bhupal Dev, R.N. Mohapatra,Phys.Rev. D 88 (2013) 033014.

[28] A. Das, P.S. Bhupal Dev, N. Okada,Phys.Lett. B735 (2014) 364. A. Das, N. Okada , Phys. Rev. D 88 (2013) 113001.

[29] Particle Data Group, J. Beringer et al., Phys. Rev. D 86 (2012) 010001.

[30] F. del Aguila and J. A. Aguilar-Saavedra, Nucl. Phys. 813 (2009) 22; F. del Aguila and J. A. Aguilar-Saavedra, Phys. Lett. B 672 (2009) 158.

[31] A.D. Martin, W.J. Stirling, R.S. Thorne, G. Watt, Eur. Phys. J. C 63 (2009) 189-285.