TeV Scale Gauged $B - L$ With Inverse Seesaw Mechanism

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We propose a modified version of the TeV scale $B - L$ extension of the standard model, where neutrino masses are generated through the inverse seesaw mechanism. We show that heavy neutrinos associated with this model can be accessible via clean signals at the LHC. The search for the extra gauge boson $Z_{B-L}$ through the decay into dileptons or two dileptons plus missing energy is studied. We also show that the $B - L$ extra Higgs can be directly probed at the LHC via a clean dilepton and missing energy signal.

The search for new physics at TeV scale is a major goal of the Large Hadron Collider (LHC). Nonvanishing neutrino masses represent a firm observational evidence of new physics beyond the standard model (SM). TeV scale Baryon minus Lepton (B-L) extension of the SM, which is based on the gauge group $SU(3)C \times SU(2)\,U(1)\,U(1)_{B-L}$, has been recently proposed \cite{1} as the simplest model beyond the SM that provides a viable and testable solution to the neutrino mass mystery of contemporary particle physics. There have been several attempts in the past to extend the gauge symmetry of the SM with $U(1)_{B-L}$, see for example Ref. \cite{2}.

In this model, three SM singlet fermions arise quite naturally due to the anomaly cancellation conditions. These three particles are accounted for right-handed neutrinos, and hence a natural explanation for the seesaw mechanism is obtained. In addition, the model also contains an extra gauge boson corresponding to $B - L$ gauge symmetry and an extra SM singlet scalar (heavy Higgs). If the scale of $B - L$ breaking is of order TeV, these new particles will lead to very interesting signatures at the LHC \cite{3, 4}. In general, the scale of $B - L$ symmetry breaking is unknown, ranging from TeV to much higher scales. However, it was proven \cite{5} that in supersymmetric framework, the scale of $B - L$ is nicely correlated with the soft supersymmetry breaking scale, which is TeV. Recently, there has been considerable interest in studying the phenomenological implications of TeV $B - L$ model at colliders \cite{4, 5, 6}.

In TeV scale $B - L$ extension of the SM, the Majorana neutrino Yukawa interaction: $\lambda_{\nu R} \chi \tilde{\nu}_R \nu_R$ induces the following masses for the right-handed neutrinos after $U(1)_{B-L}$ symmetry breaking: $M_{\nu_R} = \lambda_{\nu R} v'$, where $v' = \langle \chi \rangle$ is the vacuum expectation value (vev) of the $B - L$ symmetry breaking. Below the Electroweak (EW) symmetry breaking, Dirac neutrino masses, $m_D = \lambda_{\nu} v$, are generated. Here $v$ is the vev of the EW symmetry breaking and $\lambda_{\nu}$ are the Dirac neutrino Yukawa couplings. Therefore, the physical light neutrino masses are given by $m^2_D/M_{\nu_{LR}}$, which can account for the measured experimental results if $\lambda_{\nu} \lesssim 10^{-6}$. Such small couplings may be considered as unnatural fine-tuning. Nevertheless, they induce new interaction terms between the heavy neutrino, weak gauge boson $W$ and $Z$, and the associated leptons. These couplings play important role in the decay of lightest heavy neutrino at the LHC \cite{5, 6}. This signal is one of the striking signatures of TeV scale $B - L$ extension of the SM.

It is very important to note that the above analysis, which led to severe constraints on the neutrino Yukawa couplings, were based on the canonical type-I seesaw mechanism. In this paper, we propose a new modification for our TeV scale $B - L$ model \cite{1}, to prohibit type I seesaw and allow another scenario for generating light neutrino masses, namely the inverse seesaw mechanism \cite{10, 11}. Our modification is based on the following: (i) The SM singlet Higgs, which breaks the $B - L$ gauge symmetry, has $B - L$ charge $= -1$. (ii) The SM singlet fermion sector includes two singlet fermions with $B - L$ charges $= \pm 2$ with opposite matter parity. In this case, we will show that small neutrino masses can be generated through the inverse seesaw mechanism, without any stringent constraints on the neutrino Yukawa couplings. Therefore, a significant enhancement of the verifiability of TeV scale $B - L$ extension of the SM is obtained.

The proposed TeV scale $B - L$ extension of the SM is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_X \times U(1)_{B-L}$, where the $U(1)_{B-L}$ is spontaneously broken by a SM singlet scalar $\chi$ with $B - L$ charge $= -1$. As in the previous model, a gauge boson $Z'_{B-L}$ and three SM singlet fermions $\nu_R$, with $B - L$ charge $= -1$ are introduced for the consistency of the model. Finally, three SM singlet fermions $S_1$ with $B - L$ charge $= -2$ and three singlet fermions $S_2$ with $B - L$ charge $= +2$ are considered to implement the inverse seesaw mechanism.

The Lagrangian of the leptonic sector in this
model is given by

\[ \mathcal{L}_{B-L} = -\frac{1}{4} F_{\mu\nu}' F'^{\mu\nu} + i \bar{D}_{\mu} \gamma^\mu L + i \bar{e}_R D^\mu e_R + \bar{\nu}_R D^\mu \gamma^\mu \nu_R + i \bar{S}_1 D^\mu \gamma^\mu S_1 + i \bar{S}_2 D^\mu \gamma^\mu S_2 + (D^\mu \phi)' (D_\mu \phi) + (D^\mu \chi)' (D_\mu \chi) - V(\phi, \chi) - \left( \lambda_s \bar{L} e_R + \lambda_l \bar{S}_1 \nu_R + \lambda_s S_2 \nu_R S_2 + h.c. \right) - \frac{1}{M_3^2} S_1^4 - \frac{1}{M_3^2} S_2^4, \tag{1} \]

where \( F_{\mu\nu}' = \partial_\mu Z_{\nu}' - \partial_\nu Z_{\mu}' \) is the field strength of the \( U(1)_{B-L} \). The covariant derivative \( D_\mu \) is generalized by adding the term \( ig' Y_{B-L} Z_{\mu}' \), where \( g' \) is the \( U(1)_{B-L} \) gauge coupling constant and \( Y_{B-L} \) is the \( B - L \) quantum numbers of involved particles. Since \( U(1)_{B-L} \) is not orthogonal to \( U(1)_Y \), a mixing term between the two field strengths is expected. However, in the basis of diagonalizing kinetic terms, one finds \( ig' Y_{B-L} \to ig Y + g'' Y_{B-L} \), where \( g'' \) is parameterizing the mixing between the neutral gauge bosons: \( Z \) and \( Z' \), which is constrained experimentally to be small. Therefore, setting \( g'' = 0 \) is an acceptable approximation. In this case, what is called minimal \( B - L \) model is obtained \[1, 3\]. Furthermore, in order to prohibit a possible large mass term \( M S_1 S_2 \) in the above Lagrangian, we assume that the SM particles, \( \nu_R, \chi \), and \( S_2 \) are even under matter parity, while \( S_1 \) is an odd particle. Finally, \( V(\phi, \chi) \) is the most general Higgs potential invariant under these symmetries and can be found in Ref. \[1\].

The non-vanishing vacuum expectation value (vev) of \( \chi : |\langle \chi \rangle| = v' / \sqrt{2} \) is assumed to be of order TeV, consistent with the result of radiative \( B - L \) symmetry breaking found in gauged \( B - L \) model with supersymmetry \[4\]. The vev of the Higgs field \( \phi : |\langle \phi^0 \rangle| = v / \sqrt{2} \) breaks the EW symmetry, i.e., \( v = 246 \text{ GeV} \). After the \( B - L \) gauge symmetry breaking, the gauge field \( Z' \) acquires the following mass:

\[ M_{Z_{B-L}}^2 = g'^2 v'^2. \]

The bound on \( B - L \) gauge boson, due to negative search at LEP II, implies that \( M_{Z_{B-L}} / g'' > 6 \text{ TeV} \). This indicates that \( v' \gtrsim O(\text{TeV}) \). If the coupling \( g'' < O(1) \), then one obtains \( m_{Z'} \gtrsim O(600) \text{ GeV} \). Now, we turn to neutrino masses in this model. As can be seen from Eq. \[1\], after \( B - L \) and EW symmetry breaking, the neutrino Yukawa interaction terms lead to the following mass terms:

\[ \mathcal{L}'_{\nu} = m_D \bar{\nu}_L \nu_R + M_N \bar{\nu}_R S_2 + h.c., \tag{2} \]

where \( m_D = \frac{1}{\sqrt{2}} \lambda_d v \) and \( M_N = \frac{1}{\sqrt{2}} \lambda_{\nu_R} v' \). From this Lagrangian, one can easily observe that although the lepton number is broken through the spontaneous \( B - L \) symmetry breaking, a remnant symmetry: \(-1)^{L+S} \) is survived, where \( L \) is the lepton number and \( S \) is the spin. After this global symmetry is broken at much lower scale, a mass term for \( S_2 \) (and possibly for \( S_1 \) as well) is generated. Therefore, the Lagrangian of neutrino masses, in the flavor basis, is given by:

\[ \mathcal{L}'_{\nu} = \mu_s S_1^2 S_2 + (m_D \bar{\nu}_L \nu_R + M_N \bar{\nu}_R S_2 + h.c.), \tag{3} \]

where \( \mu_s = \frac{v''}{4 M_3} \sim 10^{-9} \). The possibility of generating small \( \mu_s \) radiatively, in general inverse seesaw model, has been discussed in Ref. \[12\].

In the basis \( \{ \nu_L, \nu_R, S_2 \} \), the \( 9 \times 9 \) neutrino mass matrix takes the form:

\[ \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & 0 & M_N \\ 0 & M_N^T & \mu_s \end{pmatrix}. \tag{4} \]

The diagonalization of this mass matrix leads to the following light and heavy neutrino masses respectively:

\[ m_{\nu_L} = m_D M_N^{-1} \mu_s (M_N^T)^{-1} m_D^T, \tag{5} \]
\[ m_{\nu_{H'}} = m_D^2 = M_N^2 + m_{D}^2. \tag{6} \]

Thus, the light neutrino mass can be of order eV, as required by the oscillation data, for a TeV scale \( M_N \), provided \( \mu_s \) is sufficiently small, \( \mu_s \ll M_N \). In this case, the Yukawa coupling \( \lambda_s \) is no longer restricted to a very small value and it can be of order one. Therefore, the possibility of testing this type of model in LHC is quite feasible.

In general, the physical neutrino states are given in terms of \( \nu_L, \nu_R, \) and \( S_2 \) as follows:

\[ \nu_1 = \nu_L + a_1 \nu_R + a_2 S_2, \tag{7} \]
\[ \nu_H = a_3 \nu_L + a_\nu_R - a_2 S_2, \tag{8} \]
\[ \nu_{H'} = a_\nu_R + a_2 S_2. \tag{9} \]

For \( m_D \approx 100 \text{ GeV}, M_N \approx 1 \text{ TeV} \) and \( \mu_s \approx 1 \text{ KeV} \), one finds that \( a_{1,2} \sim m_D / (M_N \sqrt{2 + 2 m_D / M_N}) \sim O(0.05), a_3 \sim m_D / M_N \sim O(0.1) \) and \( a_\nu \sim \sin \pi / 4 \). Therefore, one of the heavy neutrinos of this model can be accessible via a clean signal at LHC, as will be discussed below.

It is worth mentioning that the light neutrinos \( \nu_1 \) have suppressed mixing (of order \( m_D \mu_s / (M_N^2 + m_{D}^2) \)) with one type of the heavy neutrinos (say \( \nu_H \)) and a large mixing (of order \( m_D / M_N \)) with the other type of heavy neutrinos (\( \nu_{H'} \)). The mixing between the heavy neutrino \( \nu_H \) and \( \nu_{H'} \) is maximal. The heavy neutrinos \( \nu_H \) and \( \nu_{H'} \) can mediate the lepton flavor processes, like \( \mu \to e \gamma \). The \( \mu \to e \gamma \) decay mediated by these heavy neutrinos have branching
where $\Gamma_\mu$ is the total decay width of $\mu$ and the loop function $I(x)$ can be found in Ref. 13. From the present experimental limit: $BR(\mu\to e\gamma)$ one finds

$$\left| (a_3)_{\mu\mu} \right|^2 \lesssim 10^{-4}. \quad (11)$$

Thus for $(a_3)_{\mu\mu} \approx 0.1$, one obtains the following constraint on the off-diagonal element $(a_3)_{12}$:

$$(a_{12}) \simeq (m_D M_N^{-1})_{12} < 10^{-3}. \quad (12)$$

It is important to notice that within the inverse seesaw mechanism, the branching ratio of $l_\alpha \to l_\beta \gamma$ are significantly enhanced compared to the results in the conventional type I seesaw model.

The LHC discovery of $Z'_{B-L}$ is considered as a smoking gun for TeV scale $B-L$ extension of the SM. In minimal $B-L$ model, it was shown that $Z' \to l^+l^-$ gives the dominant decay channel with $BR(Z' \to l^+l^-) \approx 20\%$. Therefore, the search for $Z'$ can be accessible via a dilepton channel for 600 GeV $\leq M_{Z'} \leq$ 2 TeV. In our new model of $B-L$ with inverse seesaw, the decay widths of $Z'$ into lightest heavy neutrinos $\nu_H$ and $\nu_{H'}$ are given by:

$$\Gamma(Z' \to \nu_H\nu_H) = \frac{(g^\nu Y^\nu_{B-L})^2}{48\pi} M_{Z'} \left( 1 - 4 \frac{m^2_{\nu_H}}{M_{Z'}^2} \right)^{3/2},$$

$$\Gamma(Z' \to \nu_H'\nu_{H'}) = \frac{(g^\nu Y^\nu_{B-L})^2}{48\pi} M_{Z'} \left( 1 - 4 \frac{m^2_{\nu_{H'}}}{M_{Z'}^2} \right)^{3/2}. \quad (13)$$

From Eqs. (8,9), the charges $Y^\nu_{B-L}$ and $Y^\nu_{B-L'}$ are given by

$$Y^\nu_{B-L} \simeq a_3^2 Y^e_{B-L} + a_3 \left( Y^\nu_{B-L} - Y^\nu_{B-L} \right) \simeq 3a^2 \simeq \frac{3}{2},$$

$$Y^\nu_{B-L'} \simeq a^2 \left( Y^\nu_{B-L} + Y^\nu_{B-L} \right) - a^2 \simeq -1/2. \quad (14)$$

Thus, for heavy $Z'_{B-L}$ ($M_{Z'_{B-L}} \gg 2 M_{\nu_H}$), the decay channel $Z'_{B-L} \to \nu_H\nu_H$ could be the dominant. In Fig 1 we present the decay branching ratios of $Z' \to f\bar{f}$ as a function of $M_{Z'}$ for $f = l^-, \nu_H, \nu_l, \nu_{H'}$, $q = u, c, d, s, b$, and $f = t$. As can be seen from this figure, the decay $Z' \to l^+l^-$ is the dominant if $M_{Z'_{B-L}} < 2 M_{\nu_H}$. However, for $M_{Z'_{B-L}} \gg 2 M_{\nu_H}$, the decay $Z'_{B-L} \to \nu_{H'}\nu_{H'}$ becomes dominant with branching ratio $\gtrsim 32\%$. Therefore, searching for $Z'_{B-L}$ can be easily accessible at the LHC via: (i) A clean dilepton signal, which can be one of the first new physics signatures to be observed at the LHC, if $Z'_{B-L}$ is lighter than twice $\nu_H$ mass. As emphasized in Ref. 8, $Z'_{B-L}$ can be discovered in this case, within a mass range [800, 1200] GeV and an integrated luminosity of 100 pb$^{-1}$. (ii) A signal of 2-dilepton plus missing energy, with a tiny SM background if $M_{Z'_{B-L}} \gg 2 M_{\nu_H}$. In this case, one considers the $Z'_{B-L}$ decay into two heavy neutrinos. This process could enhance the $\nu_{H'}$ production cross section, due to the resonant contribution from $Z'_{B-L}$ exchange in the $s$-channel. Then, the $\nu_{H'}$ mainly decays through the $W$ gauge boson to lepton and neutrino, as shown in Fig 2. As explained in Ref. 8, these decays are very clean with few hard lepton, therefore they are distinctive LHC signals with nearly free background. Note that in this model, the coupling of $\nu_{H'}Wl$ is of order 0.05$g_2$, which is not very suppressed as in the minimal $B-L$ model. Therefore, the decay width of $\nu_{H'} \to W^+l^-$ is not very small, and hence $\nu_{H'}$ is no longer a long-lived particle. This could be a distinguish difference between the two $B-L$ scenarios 12.

After the breakdown of the $B-L$ and EW symmetry, mixing between $\phi$ and $\chi$ is generated. The mixing between the neutral scalar components of Higgs multiplets, $\phi^0$ and $\chi^0$, leads to the following mass eigenstates: SM-like Higgs boson $H$ and heavy Higgs boson $H'$:

$$\begin{pmatrix} H \\ H' \end{pmatrix} = \begin{pmatrix} \cos \theta & - \sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \phi^0 \\ \chi^0 \end{pmatrix},$$

(16)

where the mixing angle $\theta$ is defined by

$$\tan 2\theta = \frac{|\lambda_3| v v'}{\lambda_1 v^2 - \lambda_2 v'^2}. \quad (17)$$
The masses of $H$ and $H'$ are given by

$$m_{H,H'}^2 = \lambda_1 v^2 + \lambda_2 v'^2 \mp \sqrt{(\lambda_1 v^2 - \lambda_2 v'^2)^2 + \lambda_3^2 v^2 v'^2}. \quad (18)$$

These expressions imply that

$$m_H \approx m_{H'} \cos \theta$$

for $m_H < 120 \text{GeV}$ and $m_{H'} > 500 \text{GeV} \implies \cos \theta > 0.9. \quad (19)$$

Therefore, the cross sections of the SM-like Higgs production cross sections and decay branching ratios are slightly changed. Also, the decay widths of $H'$ into SM fermions are suppressed by $\sin^2 \theta$ factor. Due to a large mixing between light and heavy neutrinos in this model, the decay channels $H' \rightarrow \nu_H H$, $H' \rightarrow \nu_H \nu_H$ and $H' \rightarrow \nu_H \nu_{H'}$ (in case of $m_{H'} > m_{\nu_H}$, $m_{H'} > 2m_{\nu_H}$, and $m_{H'} > 2m_{\nu_{H'}}$, respectively) are relevant and may lead to important effects. The decay widths of these channels are given by

$$\Gamma(H' \rightarrow \nu_H \nu_H) = \frac{|\lambda_S a_2|^2}{32 \pi} m_{H'} \cos^2 \theta \left[1 - \frac{m_{\nu_H}^2}{m_{H'}^2}\right]^2, \quad (20)$$

and

$$\Gamma(H' \rightarrow \nu_H \nu_{H'}) \approx \Gamma(H' \rightarrow \nu_H \nu_H) \times \Gamma(H' \rightarrow \nu_{H'} \nu_{H'}) \approx \frac{|\lambda_S|^2}{64 \pi} m_{H'} \cos^2 \theta \left[1 - \frac{4m_{\nu_{H'}}^2}{m_{H'}^2}\right]^{-3/2}. \quad (21)$$

where $a_2$ is the mixing between light and heavy neutrinos as defined in Eq. (7), which is of order $0.04$. Thus, for $m_{H'} \sim 1 \text{ TeV}$, the decay width $\Gamma(H' \rightarrow \nu_H \nu_H) \sim 10^{-3}$. This should be compared with the dominant decay channel: $H' \rightarrow WW$, which has an order one decay width:

$$\Gamma(H' \rightarrow W^+ W^-) = \frac{M_{H'}^3}{16 \pi v^2} \sin^2 \theta \left[1 - \frac{4m_{W}^2}{m_{H'}^2}\right]^{3/2}. \quad (22)$$

The decay branching ratios of $H'$ into $W^+ W^-$, $ZZ$, $\nu_H \nu_H$, $\nu_{H'} \nu_{H'}$, $t\bar{t}$ and $b\bar{b}$ are shown in Fig. 3 as function of $m_{H'}$. From this figure, it is clear that the decay of $H'$ is dominated by the same channel of the SM-like Higgs. Therefore, these decay channels are experimentally challenged, due to a large background from the SM Higgs decays and can not be considered for probing $H'$ at the LHC. Furthermore, the $H'$ decay into two heavy neutrinos gives the same signal of two dileptons and missing energy as in $Z'$ decay, but with a smaller cross section. Therefore, the $H'$ production and decay via $H' \rightarrow \nu_H \nu_H \rightarrow l^+ l^- + \text{missing energy}$, as shown in Fig. 4, remains as a distinctive signal at the LHC that is nearly background free.

The total cross section of this process: $\sigma_{2l} = \sigma(pp \rightarrow H' \rightarrow \nu_H \nu_H \rightarrow l^+ l^- + \text{missing energy})$ can be written as

$$\sigma_{2l} \simeq \sigma(pp \rightarrow \tilde{\nu} \nu_H) \times BR(\nu_H \rightarrow l^- W^+) \times BR(W^+ \rightarrow l^+ \nu_l), \quad (23)$$

where $BR(W^+ \rightarrow l^+ \nu_l) \sim 0.1$ and $BR(\nu_H \rightarrow l^- W^+) \sim \mathcal{O}(1)$, since $\nu_H \rightarrow l^- W^+$ is the dominant decay channel for the heavy neutrino to the SM particles. Finally the cross section $\sigma(pp \rightarrow H' \rightarrow \nu_H \nu_H)$ can be approximated as $\sigma(pp \rightarrow H') \times BR(H' \rightarrow \nu_H \nu_{H'})$, where the $H'$ production is dominated by gluon-gluon fusion mechanism as shown in Fig. 4. In this case, $\sigma(pp \rightarrow H') \sim \mathcal{O}(0.01)$ as emphasized in Ref. 3. Also from Fig. 3 one can notice that $BR(H' \rightarrow \nu_H \nu_{H'}) \sim 10^{-3}$. Therefore, $\sigma(pp \rightarrow H' \rightarrow \nu_{H'} \nu_{H'}) \sim 10^{-5}$. In this case, the total cross section of the two dilepton signal, which provide indisputable
evidence for probing the $B-L$ extra Higgs $H'$, is give by

$$\sigma_{2l} = \sigma(pp \to H' \to l^+l^- + \text{missing energy}) \approx 10^{-7}\text{GeV}^{-2} \approx O(100)\text{pb}. \quad (24)$$

For this value of cross section, the dilepton and missing energy signal can be probed at the LHC as a a clear hint for $B-L$ extra Higgs.

It is worth mentioning that if $m_{H'} > 2m_{\nu_H}$, then the decay width $\Gamma(H' \to \nu_H\nu_H)$ becomes relevant and may be dominant. However, as mentioned above, this process leads to a signals of two dileptons with missing energy similar to the decay of $Z' \to \nu_H\nu_H$ but with a smaller cross section. Therefore, this channel is not the best for probing $H'$ at the LHC.

Finally, let us note that the above mentioned two dilptons and missing energy ($4l + E_T$) and dilepton plus missing energy ($2l + E_T$) final states are mediated by the heavy neutrinos $\nu_H$, therefore they are also clean signatures for probing $\nu_H$ at the LHC.

In conclusion, we have constructed a modified version of minimal TeV scale $B-L$ extension of the SM. In this model, the neutrino masses are generated through the inverse seesaw mechanism therefore, the neutrino Yukawa coupling is no longer constrained to be less than $10^{-6}$. Thus, the heavy neutrinos associated with this model can be quite feasible at the LHC. We have discussed the main phenomenological features of this class of models. We showed that searching for the $Z'_{B-L}$ and heavy neutrinos is accessible via $4l + E_T$ final state, while searching for the extra Higgs and also heavy neutrino can be accessible through $2l + E_T$ final state. These final states are very clean signals at LHC, with negligibly small SM background.

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