# Heavy Triplet Leptons and New Gauge Boson

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**Abstract**

A heavy triplet of leptons $(\Sigma^+, \Sigma^0, \Sigma^-)_R$ per family is proposed as the possible anchor of a small seesaw neutrino mass. A new U(1) gauge symmetry is then also possible, and the associated gauge boson $X$ may be discovered at or below the TeV scale. We discuss the phenomenology of this proposal, with and without possible constraints from the NuTeV and atomic parity violation experiments, which appear to show small discrepancies from the predictions of the standard model.

# 1  Introduction

To obtain nonzero neutrino masses so as to explain the observed atmospheric [1] and solar [2] neutrino oscillations, the minimal standard model of particle interactions is often extended to include three neutral fermion singlets, usually referred to as right-handed singlet neutrinos. If they have large Majorana masses, then the famous seesaw mechanism [3] allows the observed neutrinos to acquire naturally small Majorana masses. On the other hand, there are other equivalent ways [4, 5] to realize this effective dimension-five operator [6] for neutrino mass. For example, if we replace each neutral fermion singlet by a triplet: [5, 7]

$$\Sigma = (\Sigma^+, \Sigma^0, \Sigma^-) \sim (1, 3, 0) \tag{1}$$

under $SU(3)_C \times SU(2)_L \times U(1)_Y$, the seesaw mechanism works just as well.

If the Majorana mass of $\Sigma$ is very large, then its effect at low energies is indistinguishable from that of the canonical seesaw. On the other hand, if it is at or below the TeV energy scale, which is a natural possibility if there exists a second Higgs doublet, as shown recently [8], then there are interesting new experimental signatures for the origin of neutrino mass. In Section 2, the phenomenology of this scenario is discussed.

It is well-known [9] that in the case of one additional right-handed singlet neutrino per family of quarks and leptons, it is possible to promote $B - L$ (baryon number – lepton number) from being a global U(1) symmetry to an U(1) gauge symmetry. Similarly a new $U(1)_X$ gauge symmetry [10] is also possible here. The model is described in Section 3.

Since the $X$ gauge boson may be at or below the TeV scale, it may be responsible for some of the possible discrepancies observed in recent experiments. In Section 4, we use it to explain the NuTeV result [11] and explore its phenomenological implications. In Section 5, we do the same but using atomic parity nonconservation [12] as a constraint. In Section 6, the Higgs sector is discussed and its difference from other proposals is noted. We then

conclude in Section 7.

## 2  Heavy Triplet Leptons

Instead of the usual singlet $N_R$, consider the addition of a fermion triplet $(\Sigma^+, \Sigma^0, \Sigma^-)_R$ per family to the particle content of the standard model. The Yukawa interaction

$$\mathcal{L}_Y = f[-\phi^- \bar{\nu}_L \Sigma_R^+ + \frac{1}{\sqrt{2}}(\bar{\phi}^0 \bar{\nu}_L + \phi^- \bar{e}_L)\Sigma_R^0 + \bar{\phi}^0 \bar{e}_L \Sigma_R^-] + h.c. \tag{2}$$

means that a seesaw neutrino mass matrix is obtained, i.e.

$$\mathcal{M}_{\nu\Sigma} = \begin{bmatrix} 0 & fv/\sqrt{2} \\ fv/\sqrt{2} & M_\Sigma \end{bmatrix}, \tag{3}$$

together with $e - \Sigma$ mixing [7], i.e.

$$\mathcal{M}_{e\Sigma} = \begin{bmatrix} m_e & fv \\ 0 & M_\Sigma \end{bmatrix}. \tag{4}$$

For $v = \langle \phi^0 \rangle = 174$ GeV as in the standard model, either $M_\Sigma$ has to be very large or $f$ very small for this to give realistic neutrino masses. However, if a symmetry exists which replaces $\Phi$ in Eq. (2) with a second Higgs doublet $\eta = (\eta^+, \eta^0)$ having a very small vacuum expectation value, then $M_\Sigma$ may be at or even below the TeV scale. Such a model has already been described [8], using $N_R$. It is straightforward to apply it here in the context of $\Sigma_R$.

The idea is very simple. Assign lepton number $L = -1$ to $\eta$ but $L = 0$ to $\Phi$ and $\Sigma$, then $M_\Sigma$ is an allowed Majorana mass, and $\Phi$ is replaced with $\eta$ in Eq. (2). Let $L$ be broken by explicit *soft* terms in the Lagrangian, i.e. $\mu_{12}^2 \Phi^\dagger \eta + h.c.$, then for $m_\eta^2 > 0$ and large,

$$\langle \eta^0 \rangle = u \simeq -\frac{\mu_{12}^2 v}{m_\eta^2}. \tag{5}$$

For $\mu_{12}^2 \sim 10$ GeV$^2$, $v \sim 10^2$ GeV, $m_\eta \sim 1$ TeV, we get $u \sim 1$ MeV and $m_\nu \simeq f^2 u^2 / 2M_\Sigma \sim 1$ eV or less if $M_\Sigma \sim 1$ TeV. Note that after the explicit breaking of $L$, a residual symmetry

is still conserved, i.e. the conventional multiplicative lepton number, where $\nu$, $e$, and $\Sigma$ are odd, but $\Phi$ and $\eta$ are even. In other words, there are no unwanted $\Delta L = 1$ interactions even though $\langle \eta^0 \rangle \neq 0$.

Whereas $\Sigma^0$ has no coupling to either the photon or the $Z$ boson (as is the case with $N_R$), $\Sigma^{\pm}$ interacts with both. Hence our proposal is more easily tested experimentally than the canonical seesaw. If the mass of $\Sigma^{\pm}$ is below that of $\eta$, the former decays only via its mixing with $e^{\pm}$. Thus we expect the decay modes

$$\Sigma_i^+ \rightarrow e_j^+ Z, \ \bar{\nu}_j W^+, \quad \Sigma_i^- \rightarrow e_j^- Z, \ \nu_j W^-, \tag{6}$$

which would map out the Yukawa coupling matrix $f_{ij}$, leading to a good determination of the neutrino mass matrix itself [8], up to an overall scale. On the other hand, if $M_\Sigma > m_\eta$, then the decays

$$\Sigma_i^+ \rightarrow e_j^+ \eta^0, \ \bar{\nu}_j \eta^+ \tag{7}$$

will also determine $f_{ij}$. The subsequent decays of $\eta^+$ and $\eta^0$ occur through their small mixings with $\phi^+$ and $\phi^0$, so they are dominated by $t\bar{b}$ and $t\bar{t}$ final states and should be easily identifiable.

Since $\Sigma$ and $\eta$ have distinctive signatures once they are produced, their discoveries are primarily controlled by the size of the signal. We have estimated their pair production cross sections at the LHC and at the Tevatron via the standard Drell-Yan mechanism. The spin and color averaged matrix element squares are given by

$$\overline{M^2}_{q\bar{q} \rightarrow \eta^+ \eta^-} = \frac{1}{3} e^4 (ut - m_\eta^4) \left[ \left( \frac{Q_q}{s} + \frac{L_q L_\eta}{s - M_Z^2} \right)^2 + \left( \frac{Q_q}{s} + \frac{R_q L_\eta}{s - M_Z^2} \right)^2 \right], \tag{8}$$

where

$$L_\eta = \frac{\frac{1}{2} - \sin^2 \theta_W}{\sin \theta_W \cos \theta_W}, \quad L_q = \frac{I_q^3 - Q_q \sin^2 \theta_W}{\sin \theta_W \cos \theta_W}, \quad R_q = \frac{-Q_q \sin^2 \theta_W}{\sin \theta_W \cos \theta_W}. \tag{9}$$

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The analogous matrix element squares for $\Sigma^\pm$ pair production are

$$\overline{M^2}_{q\bar{q}\to\Sigma^+\Sigma^-} = \frac{1}{3}e^4\left[(t-M_\Sigma^2)^2\left(\frac{Q_q}{s}+\frac{L_qR_\Sigma}{s-M_Z^2}\right)^2+(u-M_\Sigma^2)^2\left(\frac{Q_q}{s}+\frac{R_qR_\Sigma}{s-M_Z^2}\right)^2\right], \quad (10)$$

where

$$R_\Sigma = \frac{1-\sin^2\theta_W}{\sin\theta_W\cos\theta_W}. \quad (11)$$

Figure 1 shows the LHC and Tevatron production cross sections of the heavy scalar pair $\eta^\pm$ and Figure 2 shows those of the heavy lepton pair $\Sigma^\pm$ as functions of their mass. We see from these figures that the final luminosity of about 300 $fb^{-1}$ at the LHC will correspond to a modest discovery limit of both $\eta^\pm$ and $\Sigma^\pm$ up to a mass of about 1 TeV.

# 3   New Gauge Boson

Consider $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$ as a possible extension of the standard model, under which each family of quarks and leptons transforms as follows:

$$(u,d)_L \sim (3,2,1/6;n_1), \quad u_R \sim (3,1,2/3;n_2), \quad d_R \sim (3,1,-1/3;n_3),$$

$$(\nu,e)_L \sim (1,2,-1/2;n_4), \quad e_R \sim (1,1,-1;n_5), \quad \Sigma_R \sim (1,3,0;n_6). \quad (12)$$

It has been shown recently [10] that $U(1)_X$ is free of all anomalies [13, 14, 15] for the following assignments:

$$n_2 = \frac{1}{4}(7n_1-3n_4), \quad n_3 = \frac{1}{4}(n_1+3n_4), \quad n_5 = \frac{1}{4}(-9n_1+5n_4), \quad n_6 = \frac{1}{4}(3n_1+n_4). \quad (13)$$

This is a remarkable and highly nontrivial result.

As shown in Ref. [10], there are 6 conditions to be satisfied for the gauging of $U(1)_X$. Three of them do not involve $n_6$ and have 2 solutions:

$$(I): \quad n_3 = 2n_1-n_2, \quad n_4 = -3n_1, \quad n_5 = -2n_1-n_2; \quad (14)$$

$$(II): \quad n_2 = \frac{1}{4}(7n_1-3n_4), \quad n_3 = \frac{1}{4}(n_1+3n_4), \quad n_5 = \frac{1}{4}(-9n_1+5n_4). \quad (15)$$

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The other 3 involve $n_6$, and they are given by

$$\frac{1}{2}(3n_1 + n_4) = \frac{1}{3}p(p+1)(2p+1)n_6, \tag{16}$$

$$6n_1 - 3n_2 - 3n_3 + 2n_4 - n_5 = (2p+1)n_6, \tag{17}$$

$$6n_1^3 - 3n_2^3 - 3n_3^3 + 2n_4^3 - n_5^3 = (2p+1)n_6^3, \tag{18}$$

where an extra right-handed fermion multiplet transforming as $(1, 2p + 1, 0; n_6)$ has been added to each family of quarks and leptons.

To find solutions to the above 3 equations, consider first $p = 0$, then Eq. (16) forces one to choose solution (I), and all other equations are satisfied with $n_1 = n_2 = n_3$ and $n_4 = n_5 = n_6$, i.e. $U(1)_{B-L}$ has been obtained. Consider now $p \neq 0$, then if solution (I) is again chosen, $n_6 = 0$ is required, which leads to $U(1)_Y$, so there is nothing new.

Now consider $p \neq 0$ and solution (II). From Eqs. (16), (17), and (18), it is easily shown that

$$\frac{4n_6}{3n_1 + n_4} = \frac{6}{p(p+1)(2p+1)} = \frac{3}{2p+1} = \left(\frac{3}{2p+1}\right)^{\frac{1}{3}}, \tag{19}$$

which clearly gives the unique solution of $p = 1$, i.e. a triplet. The fact that such a solution even exists (and for an integer value of $p$) for the above overconstrained set of conditions is certainly not a "trivial" or even "expected" result.

The $U(1)_X$ charges of the possible Higgs doublets are:

$$n_1 - n_3 = n_2 - n_1 = n_6 - n_4 = \frac{3}{4}(n_1 - n_4), \quad n_4 - n_5 = \frac{1}{4}(9n_1 - n_4), \tag{20}$$

which means that two distinct Higgs doublets are sufficient for all possible Dirac fermion masses in this model. If $n_4 = -3n_1$ is chosen, then again $U(1)_X$ will be proportional to $U(1)_Y$. However, for $n_4 \neq -3n_1$, a new class of models is now possible with $U(1)_X$ as a genuinely new gauge symmetry.

# 4 Scenario A: Neutrino-Quark Scattering

Consider $\nu q$ and $\bar{\nu} q$ deep inelastic scattering. It has recently been reported [11] by the NuTeV Collaboration that their measurement of the effective $\sin^2 \theta_W$, i.e. $0.2277 \pm 0.0013 \pm 0.0009$, is about $3\sigma$ away from the standard-model prediction of $0.2227 \pm 0.00037$. In this model, the $X$ gauge boson also contributes with

$$
\begin{aligned}
J_X^\mu &= n_1 \bar{u} \gamma^\mu \left( \frac{1 - \gamma_5}{2} \right) u + n_1 \bar{d} \gamma^\mu \left( \frac{1 - \gamma_5}{2} \right) d \\
&+ n_2 \bar{u} \gamma^\mu \left( \frac{1 + \gamma_5}{2} \right) u + n_3 \bar{d} \gamma^\mu \left( \frac{1 + \gamma_5}{2} \right) d + n_4 \bar{\nu} \gamma^\mu \left( \frac{1 - \gamma_5}{2} \right) \nu.
\end{aligned} \tag{21}
$$

Assuming very small $X - Z$ mixing ($|\sin \theta| << 1$), the effective neutrino-quark interactions are then given by

$$
\mathcal{H}_{int} = \frac{G_F}{\sqrt{2}} \bar{\nu} \gamma^\mu (1 - \gamma_5) \nu [\epsilon_L^q \bar{q} \gamma_\mu (1 - \gamma_5) q + \epsilon_R^q \bar{q} \gamma_\mu (1 + \gamma_5) q], \tag{22}
$$

where

$$
\epsilon_L^u = (1 - \xi) \left( \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right) + n_1 \zeta, \tag{23}
$$

$$
\epsilon_L^d = (1 - \xi) \left( -\frac{1}{2} + \frac{1}{3} \sin^2 \theta_W \right) + n_1 \zeta, \tag{24}
$$

$$
\epsilon_R^u = (1 - \xi) \left( -\frac{2}{3} \sin^2 \theta_W \right) + n_2 \zeta, \tag{25}
$$

$$
\epsilon_R^d = (1 - \xi) \left( \frac{1}{3} \sin^2 \theta_W \right) + n_3 \zeta, \tag{26}
$$

with

$$
\xi = 2 n_4 \sin \theta \left( 1 - \frac{M_Z^2}{M_X^2} \right) \frac{g_X}{g_Z}, \tag{27}
$$

$$
\zeta = -\sin \theta \left( 1 - \frac{M_Z^2}{M_X^2} \right) \frac{g_X}{g_Z} + 2 n_4 \left( \frac{M_Z^2}{M_X^2} \right) \frac{g_X^2}{g_Z^2}. \tag{28}
$$

The parameter $\xi$ is constrained by data at the $Z$ resonance to be very small. Using the general analysis of $Z - X$ mixing [16], we find

$$
\xi = \frac{2 s^2 c^2}{c^2 - s^2} \Delta \epsilon_1 + \frac{(c^2 - s^2)^2}{2 c^2} \Delta \epsilon_2 + \frac{s^2}{c^2} \frac{(-1 - 2 s^2 + 4 s^4)}{c^2 - s^2} \Delta \epsilon_3
$$

$$= 0.624\Delta\epsilon_1 + 0.198\Delta\epsilon_2 - 0.644\Delta\epsilon_3, \tag{29}$$

where $s \equiv \sin\theta_W$ and $c \equiv \cos\theta_W$. Given that $|\Delta\epsilon_i|$ is of order 0.001, $\xi$ is too small to make much difference in the above [17]. We thus assume $\xi = 0$ ($\sin\theta = 0$) for our subsequent discussion.

To account for the NuTeV result, i.e.

$$(g_L^{eff})^2 = (\epsilon_L^u)^2 + (\epsilon_L^d)^2 = 0.3005 \pm 0.0014, \tag{30}$$

$$(g_R^{eff})^2 = (\epsilon_R^u)^2 + (\epsilon_R^d)^2 = 0.0310 \pm 0.0011, \tag{31}$$

against the standard-model prediction, i.e.

$$(g_L^{eff})^2_{SM} = 0.3042, \quad (g_R^{eff})^2_{SM} = 0.0301, \tag{32}$$

consider the following specific model as an illustration:

$$n_1 = 1, \quad n_2 = \frac{3}{4}, \quad n_3 = \frac{5}{4}, \quad n_4 = \frac{4}{3}, \quad n_5 = -\frac{7}{12}, \quad n_6 = \frac{13}{12}. \tag{33}$$

Then

$$\Delta(g_L^{eff})^2 = -\frac{2}{3}\sin^2\theta_W\zeta + 2\zeta^2, \tag{34}$$

$$\Delta(g_R^{eff})^2 = -\frac{1}{6}\sin^2\theta_W\zeta + \frac{17}{8}\zeta^2. \tag{35}$$

To fit the experimental values, we need a negative $\Delta(g_L^{eff})^2$. From Eq. (34) we see that it reaches its maximum value at

$$\zeta = \frac{1}{6}\sin^2\theta_W, \tag{36}$$

for which

$$\Delta(g_L^{eff})^2 = -\frac{1}{18}\sin^4\theta_W = -0.0028, \tag{37}$$

$$\Delta(g_R^{eff})^2 = \frac{1}{32}\sin^4\theta_W = +0.0016, \tag{38}$$

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in very good agreement with the experimental values of $-0.0037 \pm 0.0014$ and $+0.0009 \pm 0.0011$ respectively.

Using Eqs. (28), (33), and (36), we find that

$$\frac{g_X^2}{M_X^2} = \frac{\sin^2\theta_W}{16}\frac{g_Z^2}{M_Z^2}. \tag{39}$$

Thus the production of the new gauge boson $X$ may be studied as a function of the single parameter $M_X$ in this scenario. We note first that if the $U(1)_X$ assignments of Eq. (33) apply to electrons as well, then Eq. (39) is in serious conflict with atomic parity violation and $e^+e^-$ cross sections. We must therefore attribute the NuTeV anomaly as being due to the muon (and perhaps also the tau) sector [17, 18] only. In the context of $U(1)_X$, this may be accomplished as follows. We change the electron's assignments under $U(1)_X$ to zero so that it does not couple to $X$ at all. To preserve the cancellation of anomalies, we add heavy fermions at the TeV energy scale, i.e.

$$(N, F)_L \sim (1, 2, -1/2; n_4), \quad (N, F)_R \sim (1, 2, -1/2; 0); \tag{40}$$

$$E_L \sim (1, 1, -1; 0), \quad E_R \sim (1, 1, -1; n_5). \tag{41}$$

These are prevented from coupling to the known leptons by a discrete symmetry to forbid terms such as $\bar{E}_L e_R$, etc. As a result, the lightest among them is stable, in analogy to the lightest supersymmetric particle of $R$-parity conserving supersymmetry.

The spin and color averaged matrix element square for the $X$ boson signal at the Tevatron and at the LHC is given by

$$\overline{M^2}_{q\bar{q}\to X\to f\bar{f}} = \frac{1}{3}\frac{g_X^4}{(s - M_X^2)^2 + M_X^2\Gamma_X^2}\left[n_{q_L}^2(u^2 n_{f_L}^2 + t^2 n_{f_R}^2) + n_{q_R}^2(u^2 n_{f_R}^2 + t^2 n_{f_L}^2)\right], \tag{42}$$

where

$$\Gamma_X = \frac{g_X^2}{24\pi}M_X(18n_1^2 + 9n_2^2 + 9n_3^2 + 4n_4^2 + 2n_5^2). \tag{43}$$

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Substituting the required value of $g_X$ from Eq. (39) and the $X$ charges ($n_i$) from Eq. (33) we see that for $M_X > 1$ TeV, $g_X > 1$ and its width becomes comparable to its mass. Figure 3 shows the total $X$ boson production cross sections at the LHC and at the Tevatron as functions of its mass. We see that the predicted signal cross sections are really large if the $X$ boson is to account for the NuTeV anomaly. It may be noted here that there is a 95% confidence-level upper limit of

$$\sigma(X)B(X \rightarrow e^+e^- \& \mu^+\mu^-) = 40 \ fb \qquad (44)$$

from the CDF experiment [19] at the Tevatron. The $X$ charges of Eq. (33) correspond to a branching fraction $B(X \rightarrow \mu^+\mu^-) = 4 - 5$ %. Thus assuming the CDF detection efficiency to be roughly similar for the $e^+e^-$ and $\mu^+\mu^-$ channels, the above constraint would imply $\sigma(X) < 1000 \ fb$ at the Tevatron. On the other hand we see from Figure 3 that $\sigma(X) > 1000 \ fb$ at the Tevatron right up to $M_X = 2$ TeV (the finite value at the kinematic boundary is due to the large width). Thus consistency with this limit will require $M_X > 2$ TeV. Figure 2a shows a very large signal cross section at the LHC up to $M_X = 3$ TeV, corresponding to $g_X \simeq 3$. It remains large at larger values of $M_X$ as well, the cutoff being provided by the perturbation theory limit on $g_X$.

# 5 Scenario B: Atomic Parity Violation

Instead of trying to accommodate the NuTeV discrepancy, we now consider the possibility of having a small effect in atomic parity violation from $U(1)_X$. Using the most recent precise atomic calculation of the 6s–7s parity violating E1 transition in cesium [20], a slight deviation from the standard-model prediction is obtained, i.e.

$$\Delta Q_W = 0.91 \pm 0.29 \pm 0.36. \qquad (45)$$

In the $U(1)_X$ model, with the definition $r \equiv n_1/n_4$, we have [21]

$$\Delta Q_W = -376C_{1u} - 422C_{1d} = \frac{3}{4} \frac{g_X^2}{M_X^2} \frac{M_Z^2}{g_Z^2}(1041r + 23)(1 - 9r). \tag{46}$$

This shows that there is only a narrow range for $r$ which yields a positive value of $\Delta Q_W$. For illustration, we choose $r = 0$, i.e.

$$n_1 = 0, \quad n_2 = -\frac{3}{4}, \quad n_3 = \frac{3}{4}, \quad n_4 = 1, \quad n_5 = \frac{5}{4}, \quad n_6 = \frac{1}{4}. \tag{47}$$

From Eqs. (45) and (46), we then find

$$0.015 < \frac{g_X^2}{M_X^2} \frac{M_Z^2}{g_Z^2} < 0.09. \tag{48}$$

Since this range of values will lead to large effects in the muon sector if $X$ couples in the same way, we assume in this scenario that $X$ couples only to electrons and quarks. By doing so, we also avoid serious constraints from $e^+e^- \to \mu^+\mu^-, \tau^+\tau^-$ measurements at LEP2 [22, 23].

We have estimated the $X$ boson production cross sections at the Tevatron and at the LHC using Eq. (42) for the lower limit of the $X$ coupling from Eq. (48), which is practically identical to that of Eq. (39). They are shown in Figure 4. They are somewhat smaller than those of Figure 3 due to the different $X$ charges. On the other hand, we expect a high $B(X \to e^+e^-) \simeq 18$ % in this case. Thus the CDF limit of Eq. (44) would again imply $M_X > 2$ TeV. Nonetheless we expect a very large signal cross section at the LHC up to a mass range of several TeV, till the value of $g_X$ is again cut off by the perturbation theory limit.

# 6   Higgs Sector

As shown in Sec. 3, $U(1)_X$ requires two distinct Higgs doublets for fermion masses, i.e. $\Phi_1 = (\phi_1^+, \phi_1^0)$ with $U(1)_X$ charge $(9n_1 - n_4)/4$ which couples to charged leptons, and $\Phi_2 = (\phi_2^+, \phi_2^0)$

with $U(1)_X$ charge $(3n_1 - 3n_4)/4$ which couples to *up* and *down* quarks as well as to $\Sigma$. [The leptonic Higgs doublet $\eta$ of Sec. 2 may also be introduced so that only it couples to $\Sigma$, while $\Phi_2$ couples only to quarks.] To break the $U(1)_X$ gauge symmetry spontaneously, we add a singlet $\chi$ with $U(1)_X$ charge $-2n_6$, so that the Yukawa term $\chi\Sigma\Sigma$ would allow $\Sigma$ to acquire a large Majorana mass at the $U(1)_X$ breaking scale. The Higgs potential of this model is then given by

$$
\begin{aligned}
V &= m_1^2 \Phi_1^\dagger \Phi_1 + m_2^2 \Phi_2^\dagger \Phi_2 + m_3^2 \chi^\dagger \chi + [f\chi^\dagger \Phi_1^\dagger \Phi_2 + h.c.] \\
&+ \frac{1}{2}\lambda_1(\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2}\lambda_2(\Phi_2^\dagger \Phi_2)^2 + \frac{1}{2}\lambda_3(\chi^\dagger\chi)^2 + \lambda_4(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\
&+ \lambda_5(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \lambda_6(\chi^\dagger\chi)(\Phi_1^\dagger\Phi_1) + \lambda_7(\chi^\dagger\chi)(\Phi_2^\dagger\Phi_2).
\end{aligned}
\tag{49}
$$

Note that the $U(1)_X$ charges allow the trilinear term $\chi^\dagger\Phi_1^\dagger\Phi_2$, without which $V$ would have 3 global U(1) symmetries, but only 2 U(1) gauge symmetries, resulting in an unwanted Goldstone boson. We have not included $\eta$ because $m_\eta^2$ is large and positive as discussed in Sec. 2. After the heavy $\chi$ [with $\langle\chi\rangle \sim 1$ TeV] has been integrated out, the reduced two-doublet Higgs potential is of the usual form. The difference from other proposals is in their Yukawa couplings, i.e. $\Phi_1$ couples to charged leptons whereas $\Phi_2$ couples to both *up* and *down* quarks.

Let us briefly discuss the distinctive phenomenological features of this two-Higgs-doublet model. While the vacuum expectation value $\langle\phi_2^0\rangle$ is required to be $\sim 100$ GeV because of $m_t$, $\langle\phi_1^0\rangle$ can be anywhere between $\sim 100$ GeV and 1–2 GeV. In terms of the ratio $\tan\beta = \langle\phi_2^0\rangle/\langle\phi_1^0\rangle$, they corespond to the limits $\tan\beta \simeq 1$ and $\tan\beta >> 1$. In the former case, the phenomenological implications are similar to those of the standard two-Higgs-doublet scenario where $\Phi_2$ couples to the *up* quarks and $\Phi_2$ couples to the charged leptons as well as to the *down* quarks. In the latter case however, there are distinctive differences between our proposal and the standard scenario, because Higgs couplings proportional to $m_b$ are now mutliplied by $\cot\beta$ instead of $\tan\beta$. Let us consider in particular the charged and the

pseudoscalar neutral Higgs bosons, $H^{\pm}$ and $A^0$, which correspond to the linear combination $\Phi_1 \sin\beta - \Phi_2 \cos\beta$. Their distinctive phenomenological features are summarized below.

i) The $H^- t$ loop contribution to the $b \to s\gamma$ decay amplitude is suppressed. It has the factor $\cot^2\beta$ instead of $\tan\beta \cot\beta = 1$ in the standard scenario. This means that the charged Higgs boson in this model can be relatively light.

ii) The $H^- \to \tau^- \bar{\nu}$ decay dominates over $H^- \to \bar{t}b$ for any charged Higgs mass.

iii) The $H^{\pm}$ production via $t\bar{b}$ fusion is no longer the main production mechanism at hadron colliders. It will instead be pair production via the Drell-Yan mechanism discussed in Section 2. The plots of Figure 1 apply equally to $H^{\pm}$ pair production in this case. Thus we expect a visible signal at the LHC up to a Higgs mass of about 1 TeV.

iv) The $A^0 \to \tau^+ \tau^-$ decay dominates over $A^0 \to b\bar{b}$ and $A^0 \to t\bar{t}$.

v) The $A^0$ production is again no longer dominated by $b\bar{b}$ or $t\bar{t}$ fusion at hadron colliders, but rather by associated production of $A^0 H^0 (H^{\pm})$ through $Z(W^{\pm})$ exchange.

There are analogous distinctions for the two physical neutral scalars $h^0$ and $H^0$. However, they depend on the additional mixing angle $\alpha$ which is not necessarily close to $\beta$.

# 7 Conclusion

In conclusion, we have elaborated on a recently proposed model [10] of neutrino mass, where a heavy triplet of leptons $(\Sigma^+, \Sigma^0, \Sigma^-)_R$ per family is added as the anchor of the seesaw mechanism in place of the canonical singlet $N_R$. Using a second "leptonic" Higgs doublet $\eta$ [8], both $M_\Sigma$ and $m_\eta$ can be at the TeV scale and be produced at the LHC. Experimental determination of their masses and decays to charged leptons will then map out the neutrino mass matrix.

The existence of the triplet $\Sigma$ allows a new U(1) gauge symmetry [10] to be defined, with the associated gauge boson $X$ at or below the TeV scale. The $U(1)_X$ charges of the usual quarks and leptons are fixed up to one free parameter. If $g_X^2/M_X^2$ is not too small compared to $g_Z^2/M_Z^2$, there will be corrections to the standard model at low energies coming from this new gauge symmetry. We consider two scenarios. (A) The recent NuTeV anomaly [11] is explained by $U(1)_X$. (B) The possible slight discrepancy [20] in atomic parity violation [12] is explained by $U(1)_X$. We find that (A) and (B) are mutually exclusive, i.e. we can accommodate one but not the other. The resulting constraint from either (A) or (B) on $g_X^2/M_X^2$ is such that the production of $X$ has to be very large at hadron colliders. Present limits at the Tevatron imply that $M_X$ is larger than about 2 TeV. On the other hand, it is quite possible that these anomalies have other explanations, in which case there is no hard constraint on $M_X$ or $g_X$. Nevertheless, for $g_X$ of order the electroweak coupling, the production of $M_X$ up to a few TeV can be reached at the LHC.

The $U(1)_X$ gauge symmetry requires two Higgs doublets of different $U(1)_X$ charges, such that $\Phi_1$ couples to charged leptons and $\Phi_2$ couples to both *up* and *down* quarks. This differs from the standard scenario and allows in particular the charged Higgs boson to be light in the case of large $\tan\beta$, resulting in a number of distinct phenomenological predictions.

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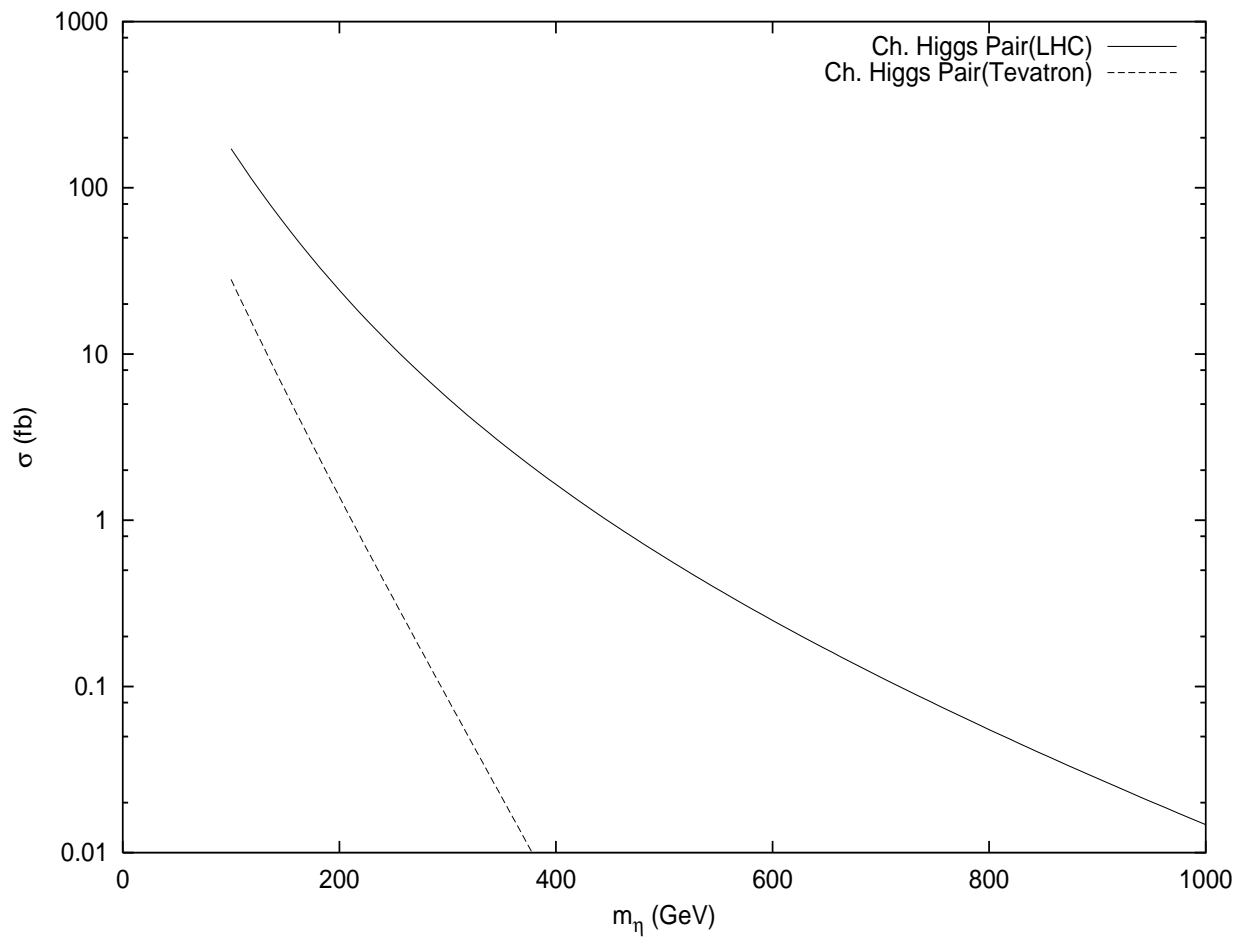
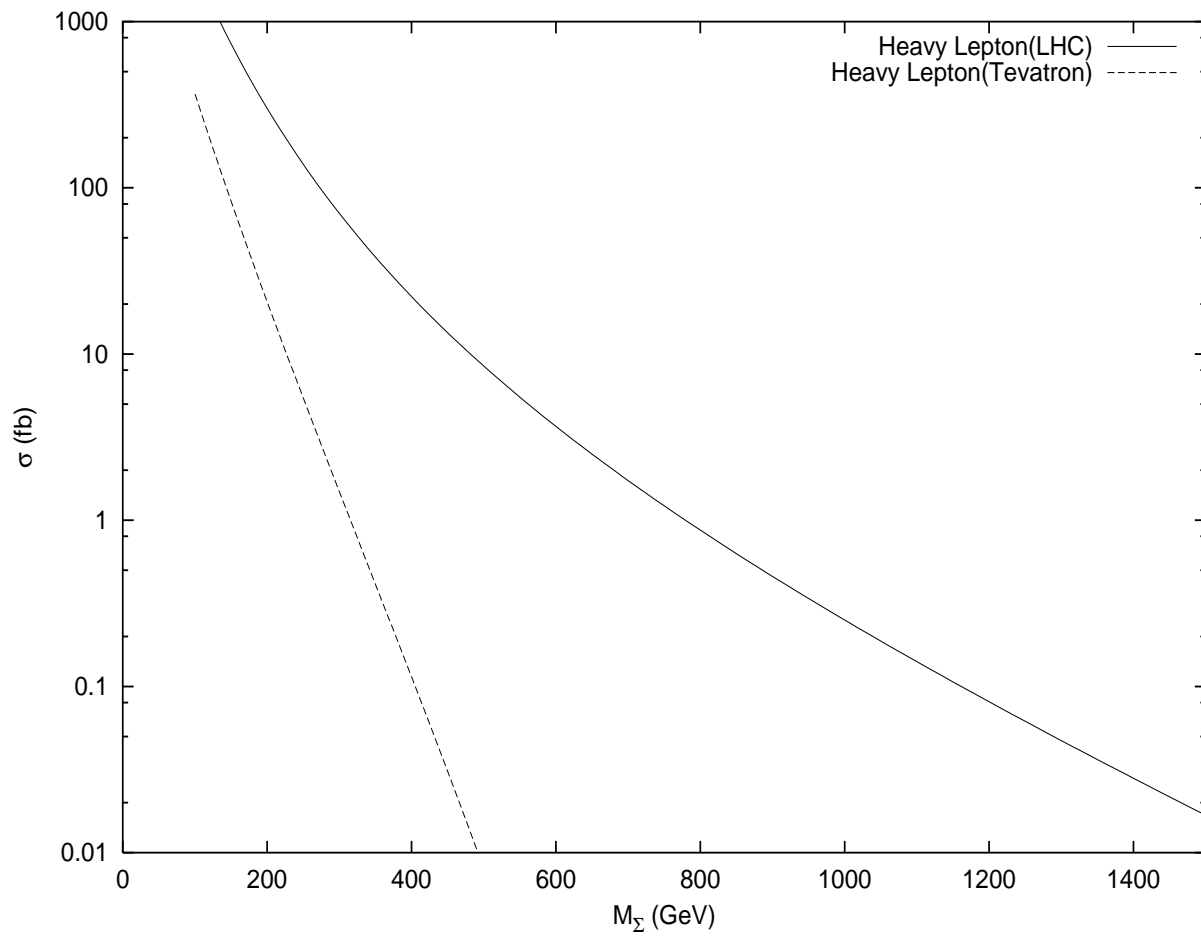Figure 1: Cross sections for $\eta^{\pm}$ pair production at LHC and Tevatron.

Figure 2: Cross sections for $\Sigma^\pm$ pair production at LHC and Tevatron.
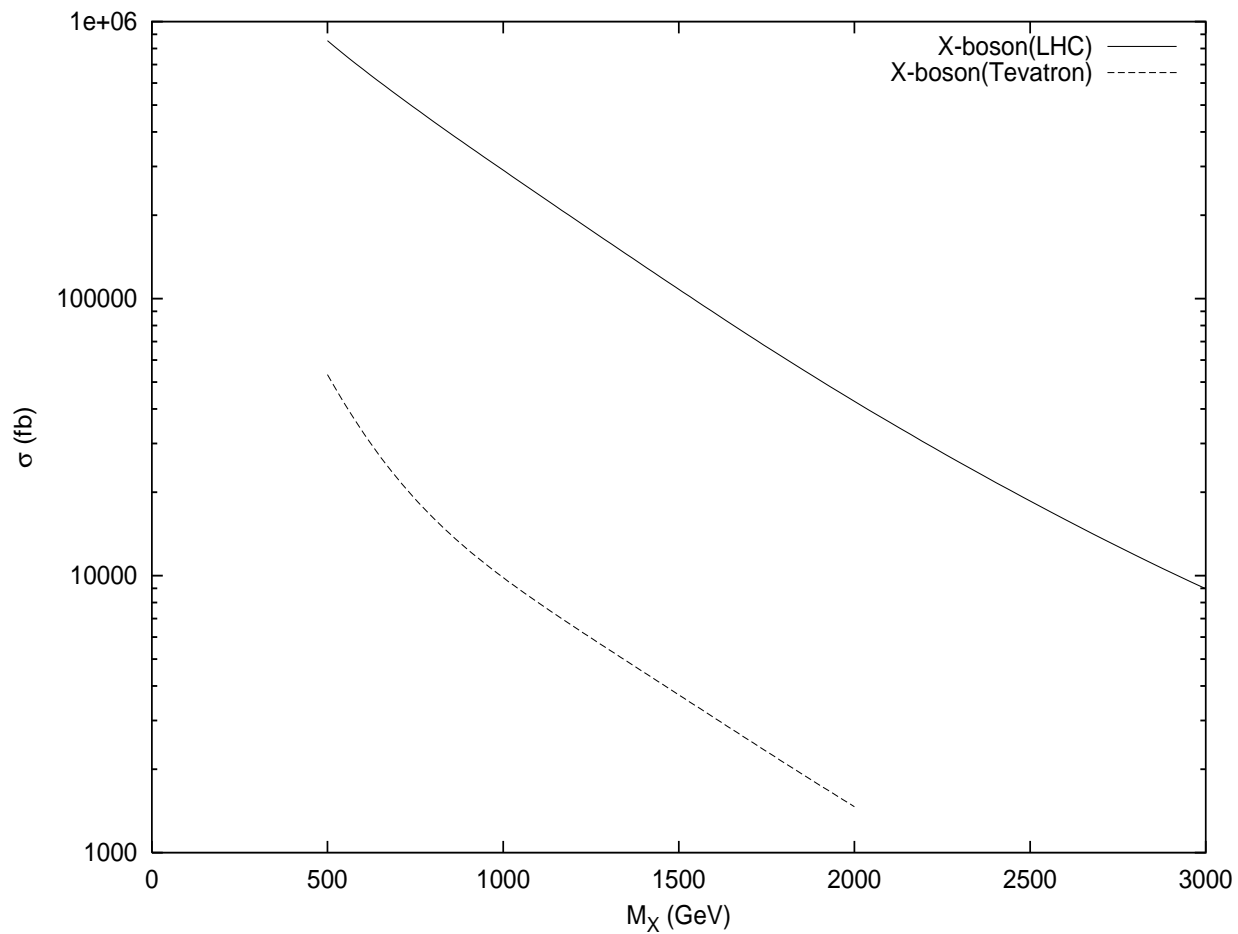
Figure 3: Cross sections for $X$ production at LHC and Tevatron in Scenario A.
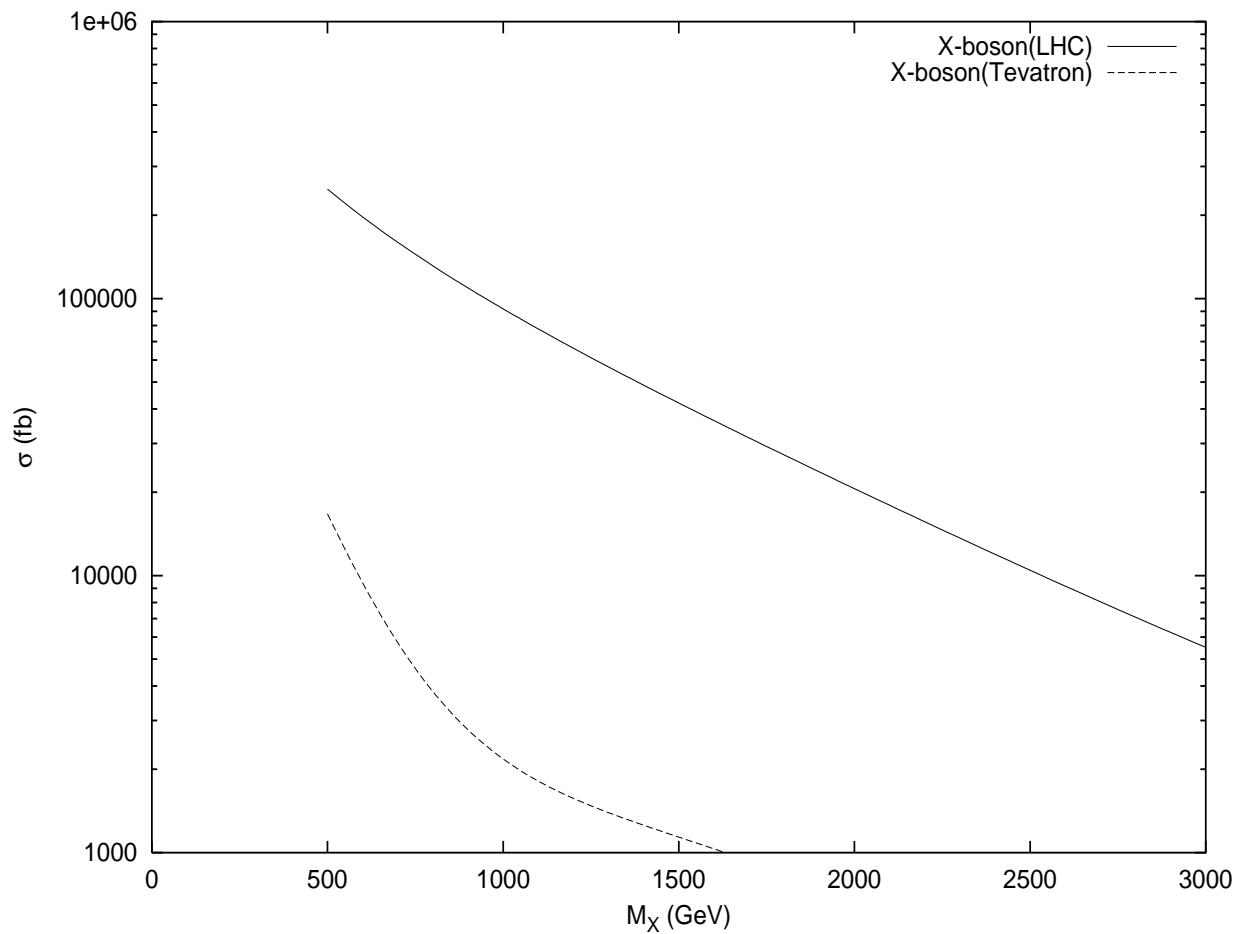
Figure 4: Cross sections for $X$ production at LHC and Tevatron in Scenario B.