Topflavor: A Separate SU(2) for the Third Family

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ABSTRACT

We consider an extended electroweak gauge group: SU(2)\textsubscript{1} × SU(2)\textsubscript{2} × U(1)\textsubscript{Y} where the first and second generation of fermions couple to SU(2)\textsubscript{1} and the third generation couples to SU(2)\textsubscript{2}. Bounds based on heavy gauge boson searches and current precision electroweak measurements are placed on the masses of the new heavy gauge bosons. In particular we find that the mass of the heavy W boson can not be less than 800 GeV. For some range of the allowed parameter space, these heavy gauge bosons produce observable signals at the Tevatron and LEP-II.

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There is no experimental data that unequivocally runs counter to the predictions of the standard SU(2)×U(1) gauge theory [1]. It is reasonable, however, to ask if a model of the electroweak interaction based on an enlarged gauge group can be constructed. In particular, we can inquire about the possibility that the electroweak symmetry breaks down from a higher symmetry at a scale that will be accessible to the next generation of colliders, the 1 TeV scale, or better yet, at present colliders such as the Tevatron or LEP II.

Perhaps the simplest extension of the standard model gauge group that one could consider is to include an extra SU(2). The model we propose is based on the gauge group SU(2)₁×SU(2)₂×U(1)ᵧ where the first and second generation of fermions couple to SU(2)₁ and the third generation couples to SU(2)₂. Other models based on the gauge group SU(2)×SU(2)×U(1) have been considered in the past; these include the left-right symmetric model [2], an un-unified model in which one SU(2) couples to quarks only while the other SU(2) couples to leptons only [3], and a model in which only one of the SU(2)’s couples to the fermions [4]. The particular model we propose is different from those above and was briefly considered by Li and Ma as an approximation to their model of generation non-universality [5]. Moreover, this particular gauge group has also been used in a non-commuting extended technicolor model [6]. Our model is somewhat analogous to the Topcolor model [7] in that the flavor sector is extended to give different weak interactions for the third family including the top quark; hence the model is named Topflavor. There is also the possibility that this new SU(2) could be used to break the Standard Model weak interaction symmetry dynamically via the Nambu-Jona-Lasinio mechanism [8].

In this work we consider the phenomenology of this extended electroweak model. We set bounds from experimental data on the masses and mixing of the new gauge bosons of this theory. We find that the new gauge bosons can be as light as 800 GeV. We then consider the possibility for detecting these new gauge bosons at present and future colliders. For some range of the SU(2)₁ and SU(2)₂ couplings, these gauge bosons give rise to detectable physics at the Tevatron and LEP-II. We then conclude with some remarks concerning flavor changing neutral current interactions. The possibilities for dynamical symmetry breaking in this model and its implications will be discussed in a future paper.

We now give a brief overview of the model. The quarks and the leptons of the first generation have the following representations under ( SU(2)₁,
SU(2), U(1)_Y:

\[
(u, d)_L \rightarrow (2, 1, 1/3), \quad u_R \rightarrow (1, 1, 4/3), \\
d_R \rightarrow (1, 1, -2/3), \quad (\nu_e, e)_L \rightarrow (2, 1, -1), \\
e_R \rightarrow (1, 1, -2).
\]  

(1)

The second generation fermions have the same representations. The fermions of the third generation have the representations:

\[
(t, b)_L \rightarrow (1, 2, 1/3), \quad t_R \rightarrow (1, 1, 4/3), \\
b_R \rightarrow (1, 1, -2/3), \quad (\nu_\tau, \tau)_L \rightarrow (1, 2, -1), \\
\tau_R \rightarrow (1, 1, -2).
\]  

(2)

With these representations for the fermions, the theory is anomaly free. The covariant derivative is

\[
D_\mu = \partial_\mu - ig'_{1/2} Y B_\mu - ig_1 T^a W^a_\mu - ig_2 \tilde{T}^b \tilde{W}^b_\mu
\]  

(3)

where the W^a belong to SU(2)_1 and the \tilde{W}^b belong to SU(2)_2.

The symmetry breaking is accomplished in two steps. First the two SU(2)'s are broken down to the SU(2)_W of the standard model (SM). Then the remaining symmetry is broken down to U(1)_em:

\[
SU(2)_1 \times SU(2)_2 \times U(1)_Y \rightarrow SU(2)_W \times U(1)_Y \rightarrow U(1)_{em}
\]  

(4)

where the electromagnetic group is generated by \( Q = T_3 + \tilde{T}_3 + Y/2 \).

The first stage in breaking the symmetry is accomplished by introducing a Higgs field \( \Phi \) that transforms as a doublet under each SU(2) with the vacuum expectation value (vev)

\[
\langle \Phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} u & 0 \\ 0 & u \end{pmatrix}.
\]  

(5)

The SM symmetry group is then broken down to U(1)_{em} through the introduction of the Higgs doublet, \( H = (1, 2, 1) \), with the vev, \( \langle H \rangle = (0, v) \). We can generate masses for the third generation of fermions with this doublet. We can generate masses for the first and second family by introducing another Higgs doublet that transforms as \( (2, 1, 1) \) although we perform the analysis here with just the one Higgs doublet.
The gauge bosons of the theory obtain mass through their interaction with the Higgs fields. The mass matrix for the neutral gauge sector is

\[
\frac{1}{2} \begin{pmatrix}
  g_1^2u^2 & -g_1g_2u^2 & 0 \\
  -g_1g_2u^2 & g_2^2(v^2 + u^2) & -g'g_2v^2 \\
  0 & -g'g_2v^2 & g'^2v^2
\end{pmatrix}
\]

where the basis is ordered as \(W, \tilde{W}, B\). This matrix can be diagonalized by means of an orthogonal matrix which we shall call \(R\):

\[
\begin{pmatrix}
  W_3 \\
  \tilde{W}_3 \\
  B
\end{pmatrix} = R^\dagger \begin{pmatrix}
  A \\
  Z_l \\
  Z_h
\end{pmatrix}
\]

where the mass eigenstate are denoted by \(A, Z_l, \) and \(Z_h\). The eigenstate \(A\) has zero mass and is identified as the photon. The couplings of our theory are related to the electric charge by

\[
g_1 = \frac{e}{\cos \phi \sin \theta_W}, \quad g_2 = \frac{e}{\sin \phi \sin \theta_W}, \quad g' = \frac{e}{\cos \theta_W}
\]

where \(\theta_W\) is the weak mixing angle and \(\phi\) is an additional mixing angle. We introduce the parameter \(\epsilon\) which is defined to be \(v^2/u^2\). Then the masses of the other two mass states are given by the equation

\[
M_{Z_i}^4 - \frac{1}{2} u^2(g_1^2 + g_2^2 + g'^2\epsilon + g'^2 \epsilon)M_{Z_i}^2 + \frac{1}{4} u^4 \epsilon(g_1^2 g'^2 + g_1^2 g'^2 + g_2^2 g'^2) = 0
\]

where \(i = l, h\). \(Z_l\) is taken as the eigenstate with the lower mass and is associated with the observed \(Z\) boson. \(Z_h\) is referred to in this paper as the "heavy \(Z\) boson". For small \(\epsilon\), the mixing matrix has the following approximate form:

\[
R = \begin{bmatrix}
  \cos \phi \sin \theta_W & \sin \phi \sin \theta_W & \cos \theta_W \\
  \cos \phi \cos \theta_W + \epsilon \cos^3 \phi \sin^2 \phi & \sin \phi \cos \theta_W - \epsilon \sin \phi \cos \phi \cos \theta_W & -\sin \theta_W \\
  -\sin \phi + \epsilon \sin \phi \cos^3 \phi & \cos \phi + \epsilon \sin^2 \phi \cos \phi \cos \theta_W & -\epsilon \tan \theta_W \sin \phi \cos^3 \phi
\end{bmatrix}
\]

The mass matrix for the charged sector is

\[
\frac{1}{2} \begin{pmatrix}
  g_1^2u^2 & -g_1g_2u^2 \\
  -g_1g_2u^2 & g_2^2(v^2 + u^2)
\end{pmatrix}
\]
where the basis is ordered as $W, \tilde{W}$. We denote by $R'$ the matrix that diagonalizes the mass matrix:

$$
\begin{pmatrix}
W \\
\tilde{W}
\end{pmatrix} = R' \begin{pmatrix}
W_l \\
W_h
\end{pmatrix}
$$

(12)

where $W_l$ and $W_h$ are the mass eigenstates whose mass are obtained by solving the equation

$$
M_{W_i}^4 - \frac{1}{2} u^2 (g_1^2 + g_2^2 + g_2^2 \epsilon) M_{W_i}^2 + \frac{1}{4} u^4 g_1^2 g_2^2 \epsilon
$$

(13)

for $i = l, h$. $W_l$ is taken as the eigenstate with the lower mass and is associated with the observed $W$ bosons. $W_h$ is referred to in this paper as the "heavy $W$ boson". For $\epsilon = 0$, $W_h$ and $Z_h$ are degenerate due to a global SU(2) symmetry. For small $\epsilon$, the mixing matrix has the approximate form

$$
R' = \begin{bmatrix}
\cos \phi + \epsilon \sin^2 \phi \cos^3 \phi & -\sin \phi + \epsilon \sin \phi \cos^4 \phi \\
\sin \phi - \epsilon \sin \phi \cos^4 \phi & \cos \phi + \epsilon \sin^2 \phi \cos^3 \phi
\end{bmatrix}.
$$

(14)

It can be easily checked that in the limit of $\epsilon = 0$, we recover the SM couplings of all the quarks and leptons to the light gauge bosons.

In this work we discuss the phenomenology in which both SU(2) interactions are perturbative. In order for this theory to be perturbative, the parameter $\phi$ can take only certain values. This range of values is delimited by $\tan \phi > 0.2$ from $\frac{g_2^2}{4\pi} < 1$ and $\tan \phi < 5.5$ from $\frac{g_2^2}{4\pi} < 1$.

We now determine what limits we can place on the two additional parameters of the theory, which we choose to be $M_{W_h}$ and $\phi$, from current experimental data. We first establish limits on the allowed values of $M_{W_h}$ and $\phi$ from precision electroweak experiments and then obtain bounds from heavy boson searches at hadronic colliders. The main result is shown in Fig. 1. The curve shown there cuts the $M_{W_h} - \tan \phi$ plane into an allowed region and a disallowed region where the upper region of the plane represents the allowed region.

We first considered bounds on the model parameters by looking at the current precision electroweak data. The quantities considered include $R_b$, $R_c$, $R_\tau$, $R_\mu$, $\sigma_h^p$ which is defined as the $Z$ boson’s cross section into hadrons at resonance, and $\frac{g_{\tau}}{g_{\mu}}$ where $g_{\tau}$ and $g_{\mu}$ are the gauge couplings of the $W_l$ to the $\tau$ and $\mu$ respectively [9]. In our model all of these processes have extra
contributions at tree level beyond the SM. Thus there is the possibility that our model’s theoretical values for these quantities could significantly deviate from the current experimental values [9], especially if the heavy Z boson mass is in the TeV scale or less.

The most restrictive of the above quantities (the only one that contributes to the bound that Fig. 1 represents) is $\frac{g_\tau}{g_\mu}$, which provides a measure of the deviation from $\mu$-$\tau$ universality: the left side of Fig. 1 was obtained by requiring that our tree level calculation for $\frac{g_\tau}{g_\mu}$ agree with the experimental value of $1.002 \pm 0.005$ [9] to within three standard deviations. In this calculation, the heavy W mass and $\tan \phi$ are input parameters. The result is that for a given value of $\tan \phi$, $\frac{g_\tau}{g_\mu}$ increases (towards unity) for increasing mass and thus sets a lower bound on the heavy W boson mass.

The right hand portion of the restriction curve is based on a heavy W boson search that was performed by the D0 Collaboration [10]. Their search involved looking for evidence of an additional W boson, which they call $W'$, decaying into an electron and its antineutrino. For various values of the $W'$ mass, they obtained an upper bound on the cross section times branching ratio for $W'$ divided by the experimental result for the same process with the usual W boson: $R \equiv \frac{\sigma_B(W' \rightarrow e\nu)}{\sigma_B(W \rightarrow e\nu)}$. The D0 collaboration considered $W'$ masses up to 800 GeV for which it established an upper bound on $R$ of $3 \times 10^{-4}$ [10].

In order to use their bounds on $R$ to set bounds on the heavy W boson mass of our model, we calculate the $W_h$ production cross section at $\sqrt{s} = 1.8$ TeV using the CTEQ distribution. Input parameters to this calculation include $\tan \phi$ and the heavy W mass. The couplings of the quarks to $W_h$ are approximated to order $\epsilon$; this is satisfactory since $\epsilon$ is small ($< 0.35$) in the region allowed by $\sigma_h^p$. The coupling of the first generation fermions to $W_h$ goes as $\tan^2 \phi$ so that the $W_h$ production cross section increases as $\phi$ increases.

Next, the branching ratio $B(W_h \rightarrow e\nu)$ is calculated. As before the input parameters to this calculation are the heavy W mass, $M_{W_h}$, and $\tan \phi$, but this time the couplings of the fermions to $W_h$ are calculated to all orders in $\epsilon$. The branching ratio for $W_h$ has additional contributions that the $W_i$ lacks from the tri-gauge boson process $W_h \rightarrow W_i Z_i$. Superficially, one might expect this process to dominate over the other contributions to the branching ratio since $\Gamma(W_h \rightarrow W_i Z_i) \sim M_h^5 / M_i^4$ for $M_h >> M_i$. However, to zeroth order
in $\epsilon$ the coupling of the above gauge bosons vanishes and, as a consequence, the contribution of this process to the branching fraction is rather small for $M_{W_i}$ up to the few TeV range. For $M_{W_i}$ around a few TeV, $B(W_i \rightarrow e\nu)$ ranges from 8 to 12.5% in the allowed range for $\tan \phi$.

Finally, we obtain our value for $R$ by taking the product of the cross section and branching ratio and then dividing by the experimental value for $\sigma_B(W \rightarrow e\nu)$ of $2.36 \pm 0.07$ nb (the value quoted by the CDF Collaboration is $2.49 \pm 0.12$) [11]. At a $W_h$ mass of 800 GeV, our values for $R$ are greater than $6 \times 10^{-4}$ for $\tan \phi > 0.8$. This is greater than the bound of $3 \times 10^{-4}$ set by the D0 Collaboration and the $R$ value gets larger with increasing $\tan \phi$. Moreover, the restriction curve from $\frac{g_{\mu}}{g}$ falls to 800 GeV around $\tan \phi = 2.1$. Therefore, we can rule out any mass for $W_h$ below 800 GeV in this model as shown in Fig. 1. This is consistent with a CDF lower bound on the mass of a new heavy W boson of 750 GeV [12]. Anticipating that the bound on $R$ for the Tevatron will improve as more data is accumulated, we show in Fig. 2 the values for $R$ that are obtained at the Tevatron for values of $M_{W_h}$ at 850, 900, 950, and 1000 GeV.

The $\rho$ parameter in this model, $\rho = M_{W_i}^2/M_{Z_i}^2 \cos^2 \theta_W$, is one to zeroth order in $\epsilon$. To all orders in $\epsilon$, the constraints from the $\rho$ parameter as well as $K^0 - \overline{K}^0$ and $B^0 - \overline{B}^0$ mixing do not restrict the allowed region further. The $\tau$ lifetime as calculated in this model is within three standard deviations of the experimental value of $\tau_{\tau} = 289.2 \pm 1.7$ fs for $\alpha_s(m_{\tau})$ in the range 0.32 to 0.38.

Next we consider the possibility for detecting the new heavy gauge bosons at the future leptonic colliders. At LEP-II the process $e^+e^- \rightarrow W_i^+W_i^-$ will be an important test of the SM prediction for the coupling of the three gauge boson coupling of the $Z$ to the $W$. In our model it is possible to have deviations from the standard model prediction for the cross section of this process due to extra terms from heavy $Z$ boson exchange as well as due to the deviation of the couplings from the SM values. Unfortunately, the $W_i$ pair production cross section at $\sqrt{s} = 200$ GeV is not significantly different from the SM value for $Z_h$ in the few TeV range. This is largely due to the fact that the coupling of $W_i$ to $Z_h$ vanishes to zeroth order in $\epsilon$ (similar to the coupling of $W_h$ to $W_i$ and $Z_i$ as mentioned above).

A process whose cross section at LEP-II can be significantly different from the SM value is $e^+e^- \rightarrow \mu^+\mu^-$ since the coupling of the electron and
Muon to the heavy Z boson goes as \( \tan \phi \). The values of the cross sections at \( \sqrt{s} = 180 \) and 200 GeV as functions of \( \tan \phi \) are shown in Fig. 3. As \( \tan \phi \) increases from 1 to 5.5, the cross section increases from close to the SM value to about twice that value. Thus, measurements of this cross section at LEP-II will either detect the effect of \( Z_h \) or eliminate a substantial part of the allowed region in Fig. 1.

We now briefly state our results for future hadronic colliders. At an upgraded Tevatron with \( \sqrt{s} = 4.0 \text{ TeV} \), the production cross section for a 1 TeV \( W_h \) increases from about 1 nb at \( \tan \phi = 1 \) to 27 nb at \( \tan \phi = 5.5 \). At the LHC with a \( \sqrt{s} = 14 \text{ TeV} \), the cross section for a 1 TeV \( W_h \) increases from 11 nb at \( \tan \phi = 1 \) to 330 nb at \( \tan \phi = 5.5 \) while, for a 3 TeV \( W_h \), the corresponding values are 0.6 and 19 nb. Thus, if the heavy gauge bosons are in these mass regions, they will be within the discovery reach of these colliders.

If the heavy \( W \) mass is in the TeV or less range, the cross section for single top production at the LHC (\( \sqrt{s} = 14 \text{ TeV} \)) through the heavy \( W \) boson resonance can be comparable to or even dominate over that for top pair production via strong interaction processes. The values we obtained for the cross section are shown in Fig. 4. We see that, for example, a 1 TeV heavy \( W \) with \( \tan \phi = 1.6 \) gives a cross section for single top production through the heavy \( W \) resonance of 1 nb which is approximately the cross section for top production through gluon fusion. An interesting signal for this single top production is two jets with a high \( p_T \) lepton. Here both jets are b-quark jets and can be tagged. This signal is different from that given by the usual \( tt \) production of the Standard Model.

Since the coupling for the third family (\( g_2 \)) differs from that of the first and second family (\( g_1 \)), this model gives rise to flavor changing neutral current (FCNC) interactions. The off-diagonal parts of the FCNC Lagrangian can be written as

\[
\mathcal{L}_{\text{neutral}}^{\text{off-diag}} = \overline{U}_L \gamma^\mu X^\dagger_L [d_l^{(u)} Z_{l,\mu} + d_h^{(u)} Z_{h,\mu}] X_L U_L \\
+ \overline{D}_L \gamma^\mu Y^\dagger_l [d_l^{(d)} Z_{l,\mu} + d_h^{(d)} Z_{h,\mu}] Y_L D_L
\]

(15)

where \( U_L \equiv (u, c, t)_L \) and \( D_L \equiv (d, s, b)_L \). Only the 3-3 elements of the four matrices \( d_l^{(u)}, d_h^{(u)}, d_l^{(d)} \) and \( d_h^{(d)} \) are nonzero and these are expressed in terms of \( g, M_{W_h} \) and \( \phi \). \( X_L \) and \( Y_L \) are defined by the following biunitary
transformations:

\[ X_L^\dagger M_u X_R = (M_u)_{\text{diag}} \]
\[ Y_L^\dagger M_d Y_R = (M_d)_{\text{diag}} . \] (16)

\(X_L\) and \(Y_L\) are related by the CKM matrix, \(K\), by \(K = X_L^\dagger Y_L\). Thus we can use the CKM matrix to eliminate \(X_l\) from eq. [15] to obtain

\[ L_{\text{neutral}}^{\text{off-diag}} = U_L \gamma^\mu K Y_L^\dagger \left[ d^{(u)}_i Z_{l,\mu} + d^{(u)}_h Z_{h,\mu} \right] Y_L K^\dagger U_L \]
\[ + D_L \gamma^\mu Y_L^\dagger \left[ d^{(d)}_i Z_{l,\mu} + d^{(d)}_h Z_{h,\mu} \right] Y_L D_L . \] (17)

Since the mass matrices are not known, we cannot calculate \(Y_L\) appearing in eq. [17]. However, using the observed hierarchy of the CKM matrix, it is easily seen that the FCNC couplings in the up and down sectors are related. For example, using \(K_{ii} \simeq 1 \gg K_{ij}\) for \(i \neq j\) and so neglecting the terms involving \(K_{ij}\), the 2-3 elements of the FCNC involving \(Z_l\) can be written as

\[ (2 - 3)_{\text{up}} \simeq \{ \bar{\sigma}_L \left[ (d^{(u)}_i)_{33} (Y_L^\dagger)_{23} (Y_L)_{33} \right] t_L + h.c. \} Z_l \]
\[ (2 - 3)_{\text{down}} \simeq \{ \bar{\sigma}_L \left[ (d^{(d)}_i)_{33} (Y_L^\dagger)_{23} (Y_L)_{33} \right] b_L + h.c. \} Z_b . \] (18)

We get similar expressions involving \(Z_h\). Thus, the FCNC interactions involving \(\bar{\sigma}\) will constrain those in the \(t\bar{c}\) sector and vice-versa.

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Figure Captions

Fig. 1. Bounds on the heavy W mass (in GeV) as a function of tan φ. The region below the curve is excluded. For tan φ < 2.1, the bounds were obtained from $g_\mu$, while the bounds for tan φ ≥ 2.1 were obtained from the limits on $R$ set by the D0 Collaboration.

Fig. 2. The values of $R$ as a function of tan φ for the Tevatron collider ($\sqrt{s} = 1.8$ TeV).

Fig. 3. The cross sections (in pb) for $e^+e^- \rightarrow \mu^+\mu^-$ as functions of tan φ for $\sqrt{s} = 180$ and 200 GeV. The SM cross sections are 3.8 and 3.0 pb respectively. For each energy the four curves correspond to $M_{Z_h} = 850, 900, 950,$ and 1,000 GeV.

Fig. 4. The cross sections (in nb) for single top production at the LHC ($\sqrt{s} = 14$ TeV). The six curves correspond to $M_{W_h} = 850, 900, 950, 1000, 2000, 3000$ GeV.