$Z\to b\bar{b}$ decay asymmetry and left-right models

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It has been pointed out recently by Chanowitz that the $Z\to b\bar{b}$ decay asymmetry poses a problem for the standard model whether or not it is genuine. If this conflict is interpreted as new physics in the b-quark couplings, it suggests a rather large right handed coupling of the b-quark to the Z boson. We show that it is possible to accommodate this result in left-right models that single out the third generation.

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I. INTRODUCTION

The precision measurements at the Z resonance continue to exhibit a deviation from the standard model in the observable A^b_{FB} by about three standard deviations [2,3]. It has been pointed out recently by Chanowitz [1] that this deviation indicates a problem whether or not it is genuine. In particular, Chanowitz argues that if the anomaly in A^b_{FB} is attributed to systematic error and dropped from the CERN e^+e^- collider LEP fits, then the indirect determination of the Higgs boson mass is in conflict with the direct limit [1].

In Ref. [4] Altarelli *et al.* approach this problem by looking for supersymmetric corrections that improve the quality of the LEP fits (including A^b_{FB}) and that improve the consistency with the direct limits on the Higgs boson mass. They find that this is possible with light sneutrinos.

The possibility of new physics affecting the Zbb coupling has also been discussed in Ref. [5]. It is known that it is not easy to explain the A^b_{FB} anomaly with new physics in the Zbb coupling mainly because the measurement of R_b is in good agreement with the standard model. However, as pointed out by Chanowitz [5], it is possible to have deviations in both the left- and right-handed couplings of the b quark to the Z boson in such a way as to change A^b_{FB} without affecting R_b.

Our starting point is the combined fit to LEP and SLAC Large Detector (SLD) measurements in terms of the left- and right-handed couplings of the b quark. These are shown in Fig. 11 of Drees [2], as well as in Ref. [3]. Subtracting the standard model values from the central value of the fit one obtains the deviations

$$\delta g_{Rb}\approx 0.02,\qquad \delta g_{Lb}\approx 0.001, \tag{1}$$

where we have flipped the sign of g_{Rb} in Refs. [2,3] to agree with the particle data book [6] definitions.

The tree-level coupling in the standard model is written as

$$L(b_L)=-\frac{g}{\cos\theta_W}\bar{b}\gamma^\mu(g_{Lb}L+g_{Rb}R)bZ_\mu, \tag{2}$$

with $L(R)=(1\mp\gamma_5)/2$. In terms of $g_V=t_{L3}-2Q\sin^2\theta_W$, $g_A=t_{L3}$ (with the parameters defined in Ref. [6]), $g_{Lb}=(g_{Vb}+g_{Ab})/2$ and $g_{Rb}=(g_{Vb}-g_{Ab})/2$. Here t_{L3} is the weak isospin which is 1/2 for up-type of quarks and $-1/2$ for down-type of quarks, and Q is the electric charge in units of e. At the tree level then

$$g_{Lb}=-\frac{1}{2}+\frac{1}{3}\sin^2\theta_W\sim -0.42,\qquad g_{Rb}=\frac{1}{3}\sin^2\theta_W\sim 0.077. \tag{3}$$

To gauge the magnitude of the required shifts, Eq. (1), it is useful to compare them with the one-loop correction in the standard model due to the heavy top quark, $\delta g_{Lb}\sim 0.003$ [5].

In view of the agreement of other low energy observables with the standard model, any new physics invoked to explain the A^b_{FB} anomaly has to affect primarily the third family, and in particular the right-handed couplings. Several scenarios in which the third generation interacts differently from the first two have been explored in the literature. Foremost amongst these is top color, where the Zbb couplings have been studied extensively in connection with R_b [7–9]. It is easy to see that while top color can easily generate a correction to the left-handed b-quark coupling of the required magnitude, it cannot generate a sufficiently large correction to the right-handed b coupling [8]. Other models considered in the literature, such as those of Refs. [10–13], single out the third family as well. However, they predominantly affect the left-handed couplings, and cannot generate the shifts required by the A^b_{FB} measurement. A possibility that may accommodate the required new physics appears in certain scenarios in which the b quark mixes with heavy quarks with unconventional charge assignments [14–17].

Alternatively, the LEP data can be attributed to new physics in the form of higher dimension operators. In this way one does not have to explain the origin of the new physics but can still use it to predict other consequences. This has been done in Ref. [18].

In this paper we explore the possibility of a left-right model that preferentially affects the third family. In Sec. II we present a model of this type and show how it can naturally accommodate the required shift in g_{Rb}. In Sec. III we explore the viability of the model in light of other existing constraints.

II. THE MODEL

The specific model to be discussed is a variation of left-right models [19,20]. The gauge group of the model is $SU(3)\times SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ with gauge couplings g_3, g_L, g_R, and g, respectively. The model differs from other left-right models in the transformation properties of the fermions.

The first two generations are chosen to have the same transformation properties as in the standard model,

$$Q_L=(3,2,1)(1/3), \quad U_R=(3,1,1)(4/3),$$
$$D_R=(3,1,1)(-2/3), \tag{4}$$
$$L_L=(1,2,1)(-1), \quad E_R=(1,1,1)(-2).$$

The numbers in the first parentheses are the $SU(3)$, $SU(2)_L$, and $SU(2)_R$ group representations, respectively, and the number in the second parentheses is the $U(1)_{B-L}$ charge.

The third generation is chosen to transform differently,

$$\begin{aligned} &Q_L(3)=(3,2,1)(1/3), \quad Q_R(3)=(3,1,2)(1/3),\\ &L_L(3)=(1,2,1)(-1), \quad L_R=(1,1,2)(-1). \end{aligned} \tag{5}$$

The above assignments are unusual compared with the conventional left-right model, but they enhance the difference between the right-handed couplings of the first two and the third generations. This model is anomaly free.

The correct symmetry breaking and mass generation of particles can be induced by the vacuum expectation values (VEVs) of three Higgs representations: $H_R=(1,1,2)(-1)$, which breaks the group down to $SU(3)\times SU(2)\times U(1)$; and the two Higgs multiplets, $H_L=(1,2,1)(-1)$ and $\phi=(1,2,2)(0)$, which break the symmetry to $SU(3)\times U(1)_{em}$. For the purpose of symmetry breaking, only one of H_L or ϕ is sufficient, but both are required to give masses to all fermions.

One may also introduce triplet Higgs multiplets, $\Delta_L=(1,3,1)(2)$ and $\Delta_R=(1,1,3)(2)$ to separate the $SU(2)_L$ and $SU(2)_R$ symmetry breaking scales and to give Majorana masses to the neutrinos. These triplets may be desirable for neutrino physics for example but they are not necessary for our present purposes.

The introduction of ϕ causes the standard model W and Z to mix with the new W_R and Z_R gauge bosons. Here W_R is the $SU(2)_R$ charged gauge boson and Z_R is a linear combination of the neutral component of the $SU(2)_R$ gauge boson W_{3R} and the $U(1)_{B-L}$ gauge boson B defined as

$$Z_R=\cos\theta_R W_{3R}-\sin\theta_R B, \tag{6}$$

where $\tan\theta_R=g/g_R$.

In the bases (W,W_R) and (Z,Z_R) for the massive gauge bosons, the mass matrices are given by

$$M_W^2=\begin{pmatrix} m_{11W}^2 & m_{12W}^2\\ m_{12W}^2 & m_{22W}^2\end{pmatrix}, \quad M_Z^2=\begin{pmatrix} m_{11Z}^2 & m_{12Z}^2\\ m_{12Z}^2 & m_{22Z}^2\end{pmatrix}, \tag{7}$$

with

$$\begin{aligned}
m_{11W}^2&=\frac{1}{2}g_L^2(|v_L|^2+2|v_{\Delta_L}|^2+|v_1|^2+|v_2|^2),\\
m_{22W}^2&=\frac{1}{2}g_R^2(|v_R|^2+2|v_{\Delta_R}|^2+|v_1|^2+|v_2|^2),\\
m_{12W}^2&=-g_Lg_R\,\mathrm{Re}(v_1v_2^*),\\
m_{11Z}^2&=\frac{1}{2}\frac{g_L^2}{\cos^2\theta_W}(|v_L|^2+4|v_{\Delta_L}|^2+|v_1|^2+|v_2|^2),\\
m_{22Z}^2&=\frac{1}{2}\frac{g_R^2}{\cos^2\theta_R}[(|v_L|^2+4|v_{\Delta_L}|^2)\sin^4\theta_R\\
&\quad+(|v_1|^2+|v_2|^2)\cos^4\theta_R+(|v_R|^2+4|v_{\Delta_R}|^2)],\\
m_{12Z}^2&=\frac{1}{4}g_Lg_R\frac{\sin\theta_R}{\cos\theta_W}[(|v_L|^2+4|v_{\Delta_L}|^2)\tan\theta_R\\
&\quad-(|v_1|^2+|v_2|^2)\cot\theta_R)],
\end{aligned} \tag{8}$$

where v_i are the VEVs of the Higgs representations $H_{L,R}$, $\Delta_{L,R}$ and ϕ.

To compare the fermion–gauge-boson couplings that result in this model with those in the standard model, we find it convenient to introduce the following definitions for gauge mixing angles:

$$\begin{aligned}
&\tan\theta_W=\frac{g_Y}{g_L}, \quad g_Y=g\cos\theta_R=g_R\sin\theta_R,\\
&\tan\theta_L=\frac{g}{g_L}, \quad \cos\theta_W=\frac{\cos\theta_L}{\sqrt{1-\sin^2\theta_L\sin^2\theta_R}},\\
&\sin\theta_W=\frac{\sin\theta_L\cos\theta_R}{\sqrt{1-\sin^2\theta_L\sin^2\theta_R}}.
\end{aligned} \tag{9}$$

After diagonalization of the mass-squared matrices, the lighter and heavier mass eigenstates (Z^1,Z^2) and (W^1,W^2) are given by

$$\begin{aligned}
\begin{pmatrix} W^1\\ W^2\end{pmatrix}&=\begin{pmatrix}\cos\xi_W & \sin\xi_W\\ -\sin\xi_W & \cos\xi_W\end{pmatrix}\begin{pmatrix} W\\ W_R\end{pmatrix},\\
\begin{pmatrix} Z^1\\ Z^2\end{pmatrix}&=\begin{pmatrix}\cos\xi_Z & \sin\xi_Z\\ -\sin\xi_Z & \cos\xi_Z\end{pmatrix}\begin{pmatrix} Z\\ Z_R\end{pmatrix},
\end{aligned} \tag{10}$$

where $\xi_{Z,W}$ are the mixing angles,

$$\tan(2\xi_{W,Z}) = \frac{2m^2_{12(W,Z)}}{m^2_{11(Z,W)} - m^2_{22(Z,W)}}. \tag{11}$$

In principle ξ_Z and ξ_W are related, and this can introduce severe constraints from processes such as $b\to s\gamma$. However, in general we find that the two can be quite different. For example, in the limit where $g \ll g_R$ we find

$$\xi_W \approx \frac{2\mathrm{Re}(v_1 v_2^\star)}{v_R^2 + 2v_{\Delta R}^2 + v_1^2 + v_2^2}\frac{\sin\theta_R}{\tan\theta_W},$$
$$\xi_Z \approx \frac{v_1^2 + v_2^2}{v_R^2 + 4v_{\Delta R}^2 + v_1^2 + v_2^2}\cos^3\theta_R \frac{\sin\theta_R}{\sin\theta_W}. \tag{12}$$

This limit is of interest because it is the one required by the A^b_{FB} data as we will see in the next section.

These results show that it is possible to have the mixing in the neutral sector be larger than the mixing in the charged sector by taking v_1 much larger (or smaller) than v_2; or by giving them a large relative phase.

The VEVs of H_L and ϕ will generate all the masses and mixings for the quarks. They also provide masses and mixings for leptons. Neutrinos in this model can also receive Majorana masses from the VEVs of Δ_L and Δ_R. If v_{Δ_R} is much larger than the electroweak scale, the right-handed neutrino will be much heavier than the left-handed neutrinos. However, there is also the possibility that the VEV of Δ_R is of the same order as the VEV of Δ_L such that all neutrinos (the three left-handed ones and the right-handed one) are light. This possibility may provide a solution to all the neutrino problems resulting from the atmospheric, solar and LSND data, should the LSND result be confirmed.

In this model there are new interactions between the massive gauge bosons and quarks. For the charged current interaction, there are both left- and right-handed interactions. In the weak eigenstate basis, the charged gauge boson W couples to all generations, but the charged gauge boson W_R only couples to the third generation. There is a similar pattern for the neutral gauge interactions. This pattern gives rise to interactions between the fermions and the lightest physical gauge bosons that can be made to resemble the standard model couplings except for the right-handed couplings of the third generation, precisely the scenario suggested by the A^b_{FB} data. In the mass eigenstate basis the quark–gauge-boson interactions are given by

$$\begin{aligned} L_W = &-\frac{g_L}{\sqrt{2}}\bar{U}_L\gamma^\mu V_{KM} D_L(\cos\xi_W W^{+1}_\mu - \sin\xi_W W^{+2}_\mu) \\ &-\frac{g_R}{\sqrt{2}}\bar{u}_{Ri}\gamma^\mu V^{u*}_{Rti}V^d_{Rbj}d_{Rj} \\ &\times(\sin\xi_W W^{+1}_\mu + \cos\xi_W W^{+2}_\mu) + \mathrm{H.c.}, \end{aligned} \tag{13}$$

where $U=(u,c,t)$ and $D=(d,s,b)$. V_{KM} is the Kobayashi-Maskawa mixing matrix and $V^{u,d}_{Rij}$ are unitary matrices which rotate the right-handed quarks u_{Ri} and d_{Ri} from the weak eigenstate basis to the mass eigenstate basis. The repeated indices i and j are summed over three generations. For the neutral sector the couplings are

$$\begin{aligned} L_Z = &-\frac{g_L}{2\cos\theta_W}\bar{q}\gamma^\mu(g_V - g_A\gamma_5)q(\cos\xi_Z Z^1_\mu - \sin\xi_Z Z^2_\mu) \\ &+\frac{g_Y}{2}\tan\theta_R\left(\frac{1}{3}\bar{q}_L\gamma^\mu q_L + \frac{4}{3}\bar{u}_{Ri}\gamma^\mu u_{Ri} - \frac{2}{3}\bar{d}_{Ri}\gamma^\mu d_{Ri}\right) \\ &\times(\sin\xi_Z Z^1_\mu + \cos\xi_Z Z^2_\mu) - \frac{g_Y}{2}(\tan\theta_R + \cot\theta_R) \\ &\times(\bar{u}_{Ri}\gamma^\mu V^{u*}_{Rti}V^u_{Rtj}u_{Rj} - \bar{d}_{Ri}\gamma^\mu V^{d*}_{Rbi}V^d_{Rbj}d_{Rj}) \\ &\times(\sin\xi_Z Z^1_\mu + \cos\xi_Z Z^2_\mu). \end{aligned} \tag{14}$$

In this expression q and q_L are summed over u,d,c,s,t,b quarks, and repeated i,j indices are summed over the three generations. The first line contains the standard model couplings to the Z in the limit $\xi_Z=0$. The first two lines also contain couplings of the two Z bosons to quarks that arise through mixing of the neutral gauge bosons.

The most interesting terms occur in the third line and are potentially large if $\cot\theta_R$ is large. In the weak interaction basis they affect only the third generation whereas in the mass eigenstate basis [as written in Eq. (14)] they also give rise to flavor-changing neutral currents. To satisfy the severe constraints that exist on flavor-changing neutral currents we must have very small V^d_{Rbd} and V^d_{Rbs} matrix elements as we discuss in the next section.

It is clear that if ξ_Z is not too small, through mixing, the b-quark coupling to the light Z boson can be very different from that of the d and s quarks due to the last term in Eq. (14). If indeed the enhancement is achieved via a large value for $\cot\theta_R$, the couplings of the first two generations will remain close to their standard model values. This illustrates how this model would solve the A^b_{FB} problem. To leading order in small ξ_Z one finds

$$\begin{aligned} \delta g_{Lb} &\approx -\frac{1}{6}\sin\theta_W\tan\theta_R\xi_Z, \\ \delta g_{Rb} &\approx \frac{1}{3}\sin\theta_W\tan\theta_R\xi_Z \\ &\quad -\frac{1}{2}\sin\theta_W(\tan\theta_R + \cot\theta_R)V^{d*}_{Rbb}V^d_{Rbb}\xi_Z. \end{aligned} \tag{15}$$

Similarly, one finds for the couplings of the top quark to the Z that, $\delta g_{Lt} = \delta g_{Lb}$, and

$$\begin{aligned} \delta g_{Rt} &\approx -\frac{2}{3}\sin\theta_W\tan\theta_R\xi_Z \\ &\quad +\frac{1}{2}\sin\theta_W(\tan\theta_R + \cot\theta_R)V^{u*}_{Rtt}V^u_{Rtt}\xi_Z. \end{aligned} \tag{16}$$

To explain the A_{FB}^b anomaly the model must be able to generate the shifts of Eq. (1). The shift required for the left-handed coupling is small and at the level of radiative corrections. We have no way of fixing all the parameters of the model at one-loop so we concentrate on the larger shift required for the right-handed coupling. Keeping only the large $\cot\theta_R$ term Eq. (1) implies that

$$\xi_Z \cot\theta_R V_{Rbb}^{d*} V_{Rbb}^{d} \sim 0.08. \tag{17}$$

We now examine this result in view of the available phenomenological constraints.

III. CONSTRAINTS

As we have pointed out, the couplings of Eq. (14) induce tree-level flavor-changing neutral currents and there are severe phenomenological constraints on these. The potentially dangerous terms in Eq. (14) are of the form

$$\frac{g_L}{2}\tan\theta_W \cot\theta_R V_{Rbi}^{q*} V_{Rbj}^{q} \overline{q}_{Ri}\gamma^\mu q_{Rj} Z_\mu^2. \tag{18}$$

The easiest way to suppress this while keeping a large Zbb right-handed coupling is by choosing the V_R^d matrix to be very close to the unit matrix. Usual constraints from $K-\overline{K}$, $D-\overline{D}$, and $B-\overline{B}$ mixing on four fermion operators, such as those generated by a tree-level exchange of Z_2, imply that off-diagonal elements in V_R^d and V_R^u are of order 10^{-4} or less [21]. Since these matrices are arbitrary, choosing them to be approximately equal to the unit matrix does not constrain other sectors of our model.

In this model there are two mechanisms that generate a large contribution to the oblique parameter T and this leads to constraints on the parameters that affect A_{FB}^b in Eq. (17). First, there is a direct contribution to T from $Z-Z_R$ mixing [22] given by

$$T=\frac{1}{\alpha}\epsilon_1=\frac{1}{\alpha}\xi_Z^2\left(\frac{M_{Z_2}^2}{M_{Z_1}^2}-1\right). \tag{19}$$

In addition there are large contributions from top (and bottom)-quark loops to the oblique corrections. Starting with the couplings in Eq. (15) these loop contributions can be obtained by extending the calculation of Ref. [27] to include the Zbb couplings as well. Starting from the effective Lagrangian

$$\begin{aligned}\mathcal{L}=&-\frac{g}{\cos\theta_W}\sum_{q=t,b}[(g_{Lq}+\delta g_{Lq})\overline{q}_L\gamma_\mu q_L\\ &+(g_{Rq}+\delta g_{Rq})\overline{q}_R\gamma_\mu q_R]Z^\mu-\frac{g}{\sqrt{2}}\{[(1+\delta\kappa_L)\overline{t}_L\gamma_\mu b_L\\ &+\delta\kappa_R\overline{t}_R\gamma_\mu b_R]W^{+\mu}+\text{H.c.}\},\end{aligned} \tag{20}$$

one finds that the leading nonanalytic contributions to the oblique parameters are [27]

$$\begin{aligned}S&=\frac{1}{6\pi}\log\left(\frac{M_{Z_2}^2}{M_Z^2}\right)[2(\delta g_{Rt}+\delta g_{Rb})-(\delta g_{Lt}+\delta g_{Lb})],\\ T&=\frac{3}{8\pi\sin^2\theta_W}\left(\frac{M_t^2}{M_W^2}\right)\log\left(\frac{M_{Z_2}^2}{M_Z^2}\right)\\ &\quad\times(2\delta\kappa_L+\delta g_{Rt}-\delta g_{Lt}),\\ U&=\frac{1}{2\pi}\log\left(\frac{M_{Z_2}^2}{M_Z^2}\right)(-4\delta\kappa_L+\delta g_{Lt}-\delta g_{Lb}).\end{aligned} \tag{21}$$

The contributions to U are seen to be small from Eq. (15). There is a potentially large contribution to S given by

$$\begin{aligned}S&=\frac{1}{6\pi}\log\left(\frac{M_{Z_2}^2}{M_Z^2}\right)\sin\theta_W\cot\theta_R\,\xi_Z V_{Rbb}^{d*}V_{Rbb}^{d}\\ &\quad\times\left(\frac{V_{Rtt}^{u*}V_{Rtt}^{u}}{V_{Rbb}^{d*}V_{Rbb}^{d}}-1\right)\\ &\sim 0.007\left(\frac{V_{Rtt}^{u*}V_{Rtt}^{u}}{V_{Rbb}^{d*}V_{Rbb}^{d}}-1\right).\end{aligned} \tag{22}$$

In the last line we have used Eq. (17) and taken $M_{Z_2}\sim 600$ GeV as a plausible upper bound (as we will see below). From Ref. [6] we know that $S=-0.03\pm0.11(-0.08)$, so new physics contributions to S are constrained to be less than 0.22 at the 2σ level. We conclude that there are no significant constraints on our model from S.

Returning to our discussion of T, we find a second large contribution,

$$T=\frac{3}{16\pi\sin\theta_W}\left(\frac{M_t^2}{M_W^2}\right)\log\left(\frac{M_{Z_2}^2}{M_Z^2}\right)\cot\theta_R\xi_Z V_{Rtt}^{u*}V_{Rtt}^{u}. \tag{23}$$

Combining Eqs. (19) and (23) and using Eq. (17) restricts the allowed $\xi_Z-M_{Z_2}$ parameter space. In line with our discussion of flavor-changing neutral currents we also require that $V_{Rtt}^{u*}V_{Rtt}^{u}/V_{Rbb}^{d*}V_{Rbb}^{d}\sim 1$. The global fit from the WWW 2001 update to Ref. [6] is

$$T=-0.02\pm0.13(+0.09). \tag{24}$$

With the particle data book definition of T this implies that new physics contributions to T are at most 0.26 at the 2σ level, and we show the resulting constraints in Fig. 1. The hatched region in the figure indicates the parameter space allowed in our model.

Additional contributions to the oblique parameters arise from the Higgs sector of our model. The model contains a standard model- (SM-) like Higgs boson from H_L which contributes in a manner similar to that of the SM Higgs boson. The remainder of the scalar sector is largely unconstrained and we shall not discuss it further in this paper.

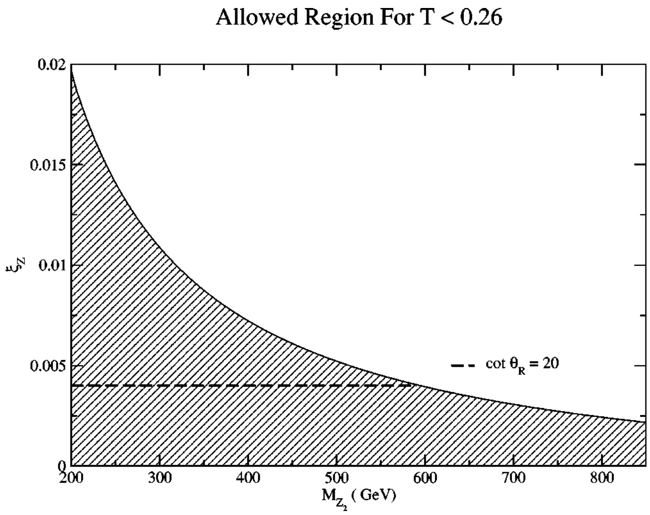

FIG. 1. Allowed region (hatched) in $\xi_Z - M_{Z_2}$ from requiring $T<0.26$, a 2σ agreement with the global fit. Below the dashed line our model becomes nonperturbative as discussed in the text.

It is instructive to discuss existing constraints on left-right (LR) models. An early comprehensive analysis of weak neutral current data [28] found that $|\xi_Z| \leqslant 0.05$ was typical for left-right models. The equivalent constraint for our model would be weaker since our Z_2 couplings to the first two generations are much weaker than in the usual LR models. Nevertheless, $|\xi_Z| \leqslant 0.05$ is satisfied in the allowed region in Fig. 1. The best direct search limits for a Z_2 boson reported in Ref. [23] come from Collider Detector at Fermilab (CDF) data [24,25] and for a LR Z_2 are of the order of 630 GeV. However, this limit assumes couplings of electroweak strength between the new Z_2 boson and the first two generations. In our model these couplings are at least ten times smaller (since they are proportional to $\tan\theta_R$ and we need $\cot\theta_R \geqslant 10$) effectively reducing the *lower bound* on the Z_2 mass to less than 100 GeV from experiments that only involve couplings to the first two generations. Reference [13] has studied the problem of placing bounds on the mass of a Z_2 that couples preferentially to the third generation. From searches for compositeness there are bounds on the scale of four fermion operators such as $\bar{q}_L \gamma_\mu q_L \bar{e}_L \gamma^\mu e_L$ on the order of 4 TeV [26]. However, when these operators are induced by the exchange of the Z_2 in our model, the smallness of its couplings to the first two generations results in a very weak lower bound (well below 100 GeV) as well.

It appears that there are no significant constraints on the mass of a Z_2 with couplings to the first two generations as weak as those in our model. For illustration we will simply assume a lower limit on the mass of our W_2 of order 100 GeV because it has not been produced in e^+e^- colliders, and assume a similar lower limit on the Z_2 mass to produce Fig. 1.

It is possible to place a theoretical constraint on the parameters of the model by requiring the coupling of the Z_2 to the third generation fermions to remain perturbative. Demanding that

$$\left(\frac{g_Y \cot\theta_R}{2}\right)^2 \leqslant 4\pi \tag{25}$$

results in $\cot\theta_R \leqslant 20$. Combined with Eq. (17), and assuming that $V^{d*}_{Rbb} V^d_{Rbb} \sim 1$, this results in a lower limit on $\xi_Z \geqslant 0.004$ shown as a dashed line in Fig. 1. As seen in Fig. 1, this also implies that $M_{Z_2} \leqslant 600$ GeV. The hatched region *below* the dashed line in Fig. 1 satisfies phenomenological constraints but implies a Z_2 coupling to the third generation which is nonperturbative.

We now turn our attention to the charged gauge-boson sector. The early bounds on W-W_R mixing from a comprehensive analysis of low energy data can be found in Ref. [29]. Depending on the model their typical result was

$$|\xi_g| \leqslant 10^{-3} \tag{26}$$

where

$$\xi_g \equiv \frac{g_R}{g_L}\xi_W = \tan\theta_W / \sin\theta_R \xi_W .$$

As with the neutral gauge-boson sector, these constraints do not apply directly to our model. The best bound on W_R couplings to third generation quarks comes from $b\rightarrow s\gamma$ as in the bound on the anomalous coupling $\delta\kappa_R$ from Eq. (20) obtained in Ref. [30]. A more careful treatment of QCD corrections can be done along the lines of Ref. [31]. The dominant contribution to $b\rightarrow s\gamma$ and the associated $b\rightarrow sg$ is from W-W_R mixing, one has

$$H_{\rm mixing} = -\frac{4G_F}{\sqrt{2}} V_{tb} V^*_{ts} [c_7^{LR} O_7 + c_8^{LR} O_8],$$

$$O_7 = \frac{e}{16\pi^2} m_b \bar{s} \sigma_{\mu\nu} R F^{\mu\nu} b, \tag{27}$$

$$O_8 = \frac{g_3}{16\pi^2} m_b \bar{s} \sigma_{\mu\nu} R G^{\mu\nu} b,$$

where $c^{LR}_{7,8}$ are the Wilson coefficients due to left-right mixing evaluated at a scale of order $O(m_W)$. In our model they are given by

$$c_7^{LR} = \xi_{eff} \frac{m_t}{m_b} \tilde{F}(x_t), \quad c_8^{LR} = \xi_{eff} \frac{m_t}{m_b} \tilde{G}(x_t), \tag{28}$$

where

$$\xi_{eff} = \frac{\tan\theta_W}{\sin\theta_R} \xi_W \left(\frac{V^u_{Rtt} V^{*d}_{Rbs}}{V^*_{ts}} + \frac{V^{*u}_{Rtt} V^d_{Rbb}}{V_{tb}} \right),$$

$$\tilde{F}(x) = \frac{-20 + 31x - 5x^2}{12(x-1)^2} + \frac{x(2-3x)}{2(x-1)^3} \ln x, \tag{29}$$

$$\tilde{G}(x) = -\frac{4+x+x^2}{4(x-1)^2} + \frac{3x}{2(x-1)^3} \ln x,$$

with $x_t = m_t^2/m_W^2$, and V_{tb}, V_{ts} the usual Cabibbo-Kobayashi-Maskawa (CKM) angles.

Running down to the scale relevant to B decays, we obtain the dominant effective Wilson coefficient for $b \to s\gamma$, c_{7eff},

$$c_{7eff} = \eta^{16/23} c_7^{LR} + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) c_8^{LR}. \tag{30}$$

Here $\eta = \alpha_s(m_W)/\alpha_s(m_b)$.

Compared with the SM top quark contribution, there is an enhancement factor m_t/m_b. Using the most recent experimental data for $b \to s\gamma$, $B(b \to s\gamma) = (3.21 \pm 0.43 \pm 0.27^{+0.18}_{-0.10}) \times 10^{-4}$ [32] we find at the 2σ level that there are two allowed ranges for ξ_{eff}. They correspond to destructive and constructive interference with the standard model amplitude, respectively, and are

$$\begin{aligned} -0.032 &< \xi_{eff} < -0.027, \\ -0.0016 &< \xi_{eff} < 0.0037. \end{aligned} \tag{31}$$

In line with our discussion of flavor-changing neutral currents we assume that the $V_R^{u,d}$ matrices are very close to the unit matrix. The largest contribution is then from

$$\xi_{eff} \sim \frac{\tan\theta_W}{\sin\theta_R} \xi_W. \tag{32}$$

With $\cot\theta_R \sim 10$ the two allowed ranges for ξ_W are

$$\begin{aligned} -0.006 &< \xi_W < -0.005, \\ -0.0003 &< \xi_W < 0.0007. \end{aligned} \tag{33}$$

For comparison, the bound on ξ_Z from T combined with the perturbative requirement resulted in $0.004 < \xi_Z < 0.02$. In the first allowed range in Eq. (33) one finds $0.25 < |\xi_W/\xi_Z| < 1.5$ making both mixing angles of the same order. On the other hand, in the second allowed range in Eq. (33) $0 < |\xi_W/\xi_Z| < 0.18$, so ξ_W is typically much smaller. An analysis of the quark mass matrices in our model reveals that it is possible to accommodate naturally a hierarchy $v_1/v_2 \sim m_t/m_b$.[1] This scenario would result in a $|\xi_W/\xi_Z|$ as low as $2m_b/m_t \sim 0.046$.

In conclusion we have shown that it is possible to accommodate the A_{FB}^b result in a model with new right-handed interactions for the third generation. The model predicts large deviations from the standard model in the right-handed couplings of the top-quark and perhaps the tau lepton. A striking feature of the model is the possible existence of a light ($M_{Z_2} < 600$ GeV) Z' boson on which there is no meaningful lower bound at present. The range allowed by $b \to s\gamma$ in Eq. (31) also indicates that this model can give rise to CP asymmetries in B decays that deviate significantly from the standard model [34].

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