The $Z \to b\bar{b}$ Decay Asymmetry and Left-Right Models

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Abstract

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 it is genuine or not [1]. 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.

PACS numbers:
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 it is genuine or not. In particular, Chanowitz argues that if the anomaly in $A^b_{FB}$ is attributed to systematic error and dropped from the LEP fits, then the indirect determination of the Higgs mass is in conflict with the direct limit [1].

In Ref. [4], Altarelli et.al. approach this problem by looking for super-symmetric 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 [3], 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 SLD measurements in terms of the left and right handed couplings of the $b$-quark. These are shown in Figure 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,

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

(1)

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

The tree-level coupling in the standard model is written as
\[ L(b_L) = \frac{g}{\cos \theta_W} \bar{b}_\gamma \mu (g_{Lb} L + g_{Rb} R) b Z_\mu, \]  

(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 tree-level then,

\[
\begin{align*}
    g_{Lb} &= -\frac{1}{2} + \frac{1}{3} \sin^2 \theta_W \sim -0.42 \\
    g_{Rb} &= \frac{1}{3} \sin^2 \theta_W \sim 0.077
\end{align*}
\]

(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 \( A_{FB} \) 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 \( R_b \) [7,8]. 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. [9], [10] and [11], single out the third family as well. However, they predominantly affect the left-handed couplings, and cannot generate the shifts required by the \( A_{FB} \) 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 [12–14].

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. [15].

In this paper we explore the possibility of a left-right model that preferentially affects the third family. In Section 2 we present a model of this type and show how it can naturally
accommodate the required shift in $g_{\text{RB}}$. In Section 3 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 [16,17]. 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,

\[ Q_L = (3,2,1)(1/3), \quad U_R = (3,1,1)(4/3), \quad D_R = (3,1,1)(-2/3), \]
\[ L_L = (1,2,1)(-1), \quad E_R = (1,1,1)(-2). \] (4)

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

The third generation is chosen to transform differently,

\[ Q_L(3) = (3,2,1)(1/3), \quad Q_R(3) = (3,1,2)(1/3), \]
\[ L_L(3) = (1,2,1)(-1), \quad L_R = (1,1,2)(-1). \] (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 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 $SU(3) \times U(1)_{\text{em}}$. For the purpose of symmetry breaking, only one of $H_L$ or $\phi$ 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. The introduction of $\phi$ 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

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

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,

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

with

$$m_{11W}^2 = \frac{1}{2} g_L^2 (|v_L|^2 + 2|v_{\Delta_L}|^2 + |v_1|^2 + |v_2|^2);$$

$$m_{22W}^2 = \frac{1}{2} g_R^2 (|v_R|^2 + 2|v_{\Delta_R}|^2 + |v_1|^2 + |v_2|^2);$$

$$m_{12W}^2 = -g_L g_R \Re(v_1 v_2^*) ;$$

$$m_{11Z}^2 = \frac{1}{2} \frac{g_L^2}{\cos^2 \theta_W} (|v_L|^2 + 4|v_{\Delta_L}|^2 + |v_1|^2 + |v_2|^2);$$

$$m_{22Z}^2 = \frac{1}{2} \frac{g_R^2}{\cos^2 \theta_R} ((|v_L|^2 + 4|v_{\Delta_L}|^2) \sin^4 \theta_R + (|v_1|^2 + |v_2|^2) \cos^4 \theta_R$$

$$+ (|v_R|^2 + 4|v_{\Delta_R}|^2));$$

$$m_{12Z}^2 = \frac{1}{4} g_L g_R \frac{\sin \theta_R}{\cos \theta_W} ((|v_L|^2 + 4|v_{\Delta_L}|^2) \tan \theta_R - (|v_1|^2 + |v_2|^2) \cot \theta_R)), \quad (8)$$

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

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,
\[
\tan \theta_W = \frac{g_Y}{g_L}, \quad g_Y = g \cos \theta_R = g_R \sin \theta_R, \quad \tan \theta_L = \frac{g}{g_L},
\]
\[
\cos \theta_W = \frac{\cos \theta_L}{\sqrt{1 - \sin^2 \theta_L \sin^2 \theta_R}}, \quad \sin \theta_W = \frac{\sin \theta_L \cos \theta_R}{\sqrt{1 - \sin^2 \theta_L \sin^2 \theta_R}}.
\]

(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
\[
\begin{pmatrix}
W^1 \\
W^2
\end{pmatrix}
= \begin{pmatrix}
\cos \xi_W & \sin \xi_W \\
-\sin \xi_W & \cos \xi_W
\end{pmatrix}
\begin{pmatrix}
W \\
W_R
\end{pmatrix},
\]
\[
\begin{pmatrix}
Z^1 \\
Z^2
\end{pmatrix}
= \begin{pmatrix}
\cos \xi_Z & \sin \xi_Z \\
-\sin \xi_Z & \cos \xi_Z
\end{pmatrix}
\begin{pmatrix}
Z \\
Z_R
\end{pmatrix},
\]

(10)

where \(\xi_{Z,W}\) are the mixing angles,
\[
\tan(2\xi_{W,Z}) = \frac{2m^2_{12(W,Z)}}{m^2_{11(Z,W)} - m^2_{22(Z,W)}}.
\]

(11)

In principle \(\xi_Z\) and \(\xi_W\) are related, and this can introduce severe constraints from processes such as \(b \rightarrow s\gamma\). However, in general we find that the two can be quite different. For example, in the limit where \(g << g_R\) we find
\[
\xi_W \approx \frac{2\text{Re}(v_1 v_2^*)}{v_R^2 + 2v_{\Delta R}^2 + v_1^2 + v_2^2} \tan \theta_W \sin \theta_R
\]
\[
\xi_Z \approx \frac{v_1^2 + v_2^2}{v_R^2 + 4v_{\Delta R}^2 + v_1^2 + v_2^2} \cos^3 \theta_R \frac{\sin \theta_R}{\sin \theta_W}.
\]

(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 \(\phi\) 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 \(\Delta_L\) and \(\Delta_R\). If \(v_{\Delta_R}\) 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 \(\Delta_R\) is of the same order as the vev of \(\Delta_L\) 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 $A_{FB}^b$ data. In the mass eigenstate basis the quark-gauge-boson interactions are given by,

$$L_W = -\frac{g_L}{\sqrt{2}} \bar{u}_L \gamma^\mu V_{K_M} D_L (\cos \xi_W W^1_\mu - \sin \xi_W W^2_\mu)$$

$$- \frac{g_R}{\sqrt{2}} \bar{u}_R \gamma^\mu V_{Ri}^u V_{Rj}^d (\sin \xi_W W^1_\mu + \cos \xi_W W^2_\mu),$$

(13)

where $U = (u, c, t)$ and $D = (d, s, b)$. $V_{K_M}$ is the Kobayashi-Maskawa mixing matrix and $V_{Ri}^u, V_{Rj}^d$ are unitary matrices which rotate the right handed quarks from the weak eigenstate basis to the mass eigenstate basis. For the neutral sector the couplings are,

$$L_Z = -\frac{g_L}{2 \cos \theta_W} \bar{q} \gamma^\mu (g_V - g_A \gamma_5) q (\cos \xi_Z Z^1_\mu - \sin \xi_Z Z^2_\mu)$$

$$+ \frac{g_Y}{2} \tan \theta_R (B - L) (\bar{q}_L \gamma^\mu q_L + \bar{q}_R \gamma^\mu q_R) (\sin \xi_Z Z^1_\mu + \cos \xi_Z Z^2_\mu)$$

$$- \frac{g_Y}{2} \cot \theta_R (\bar{t}_R \gamma^\mu t_R - \bar{b}_R \gamma^\mu b_R) (\sin \xi_Z Z^1_\mu + \cos \xi_Z Z^2_\mu).$$

(14)

The first two terms apply to all the quarks, $q = u, d, c, s, t, b$, with $B - L$ the respective $U(1)_{B-L}$ quantum number. The last term affects only the third generation and can be large if $\cot \theta_R$ is large.

It is clear that if $\xi_Z$ 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 the above expression. If 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 $A_{FB}^b$ problem. To leading order in a small $\xi_Z$ expansion one finds,

\begin{align}
\delta g_{Lb} &\approx -\frac{1}{6} \sin \theta_W \tan \theta_R \xi_Z \\
\delta g_{Rb} &\approx -\frac{1}{6} \sin \theta_W \tan \theta_R \xi_Z - \frac{1}{2} \sin \theta_W \cot \theta_R \sin \xi_Z.
\end{align}

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

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 quite 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 much larger shift required for the right-handed coupling, it implies that,

$$\xi_Z \cot \theta_R \sim 0.08.$$  

(16)

We now examine this result in view of the known constraints on left-right models.

III. CONSTRAINTS

As we have pointed out, it is possible to arrange the parameters in the model in such a way that the mixing in the neutral gauge boson sector is much larger than that in the charged gauge boson sector. For this reason we will treat separately the constraints on $Z - Z_R$ mixing and on $W - W_R$ mixing.

An early comprehensive analysis of weak neutral current data [18] found that $|\xi_Z| \leq 0.05$ was typical for left-right models. Since then, the most important new constraint arises from LEP data [3] and corresponds to the $Z - Z_R$ mixing contribution to the parameter $T$ [19]

$$\epsilon_1 = \alpha T = \xi_Z^2 \left( \frac{M_{Z_2}^2}{M_{Z_1}^2} - 1 \right) \leq 0.0011$$

(17)

where the last inequality follows from allowing the new contribution to be as large as the error in the Particle Data Book fit [3].
The best direct search limits for a $Z_2$ boson reported in Ref. [20] came from CDF data [21] and for a LR $Z_2$ are of the order of 450 GeV. If we use this number in Eq. [17] we find that $|\xi_Z| \leq 0.007$. Eq. [16] then tells us that

$$\cot \theta_R \geq 11.4 \text{ or } g_R \geq 6.6g.$$ \hspace{1cm} \text{(18)}$$

In addition to $Z-Z_R$ mixing, there are other sources of new contributions to the oblique parameters in this model. First, there are the scalar loop contributions; we will not be able to say much about them because the Higgs sector of the model is largely unconstrained. Second, there are the top (and bottom)-quark loops which may differ significantly from their standard model values. The easiest way to quantify this contribution is by recasting the couplings, Eq. [15], in terms of the general analysis of anomalous couplings as in Ref. [22],

$$\mathcal{L} = -\frac{g}{2 \cos \theta_W} \sum_{q=t,b} \left( (L_q + \delta L_q) \bar{q} L \gamma_\mu q_L + (R_q + \delta R_q) \bar{q} R \gamma_\mu q_R \right) Z^\mu$$

$$- \frac{g}{\sqrt{2}} \left[ (1 + \delta \kappa_L) \bar{t} L \gamma_\mu b_L + \delta \kappa_R \bar{t} R \gamma_\mu b_R \right] W^{+\mu} + \text{h.c.}. \hspace{1cm} \text{(19)}$$

The logarithmic terms found in Ref. [22] will serve to gauge the size of these contributions,

$$S = \frac{1}{3\pi} \log\left( \frac{\mu^2}{M_Z^2} \right) \left( 2(\delta R_t + \delta R_b) - (\delta L_t + \delta L_b) \right)$$

$$T = \frac{3}{4\pi \sin^2 \theta_W} \left( \frac{M_t^2}{M_W^2} \right) \log\left( \frac{\mu^2}{M_Z^2} \right) \left( \delta \kappa_L + \delta R_t - \delta L_t \right)$$

$$U = \frac{1}{\pi} \log\left( \frac{\mu^2}{M_Z^2} \right) \left( -2\delta \kappa_L + \delta L_t - \delta L_b \right) \hspace{1cm} \text{(20)}$$

There is no large contribution to $S$ in this model because $\delta R_t = -\delta R_b + \mathcal{O}(\tan \theta_R)$. The largest effect occurs in the contribution to $T$ and is given by, to $T$,

$$T = \frac{3}{4\pi \sin^2 \theta_W} \left( \frac{M_t^2}{M_W^2} \right) \log\left( \frac{M_Z^2}{M_{Z_2}^2} \right) g_R.$$

If we require once more that this be smaller than the error quoted by the Particle Data Book, $|T| \leq 0.14$, [8] and we use $M_{Z_2} = 450$ GeV [20], we find

$$\xi_Z \cot \theta_R \leq 0.07. \hspace{1cm} \text{(22)}$$
which is compatible with what is required to solve the $A_{FB}^b$ problem, Eq. [10].

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. [23]. Depending on the model their typical value was,

$$|\xi_g| \leq 10^{-3}$$

(23)

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

A similar constraint arises, specifically for the model we discuss, from the new contributions to $b \to s\gamma$. The dominant contribution to $b \to s\gamma$ and the associated $b \to sg$ is from $W - W_R$ mixing [24]. One has

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

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

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

(24)

where $c_7^{LR}, c_8^{LR}$ are the Wilson Coefficients due to Left-Right mixing evaluated at a scale of order $O(m_W)$. They are given by

$$c_7^{LR} = \xi_g \frac{m_t}{m_b} \tilde{F}(x_t) \frac{V_{tb}^d V_{ts}^*}{V_{tb}},$$

$$c_8^{LR} = \xi_g \frac{m_t}{m_b} \tilde{G}(x_t) \frac{V_{tb}^d V_{ts}^*}{V_{tb}},$$

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

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

(25)

where $x_t = m_t^2/m_W^2$.

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

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

(26)

Here $\eta = \alpha_s(m_W^1)/\alpha_s(m_b)$.  

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Compared with the SM top quark contribution, there is an enhancement factor $m_t/m_b$. The parameters $V_{Rbb}$ and $V_{Rtt}$ are close to one if the quark mixings are small. Using experimental data $(3.21 \pm 0.43 \pm 0.27^{+0.18}_{-0.10}) \times 10^{-4}$ [24], the parameter $\xi_g$ is constrained to be

$$|\xi^W| \leq 10^{-3} \frac{\sin \theta_R}{\tan \theta_W}$$

(27)

With this constraint, the $W - W_R$ mixing contribution to $T$ is small. Using Eq. 18, this implies that $|\xi^W| \leq 2 \times 10^{-4}$, which is about 35 times stronger than the bound on $\xi_Z$. As mentioned earlier, this hierarchy is allowed in the model. One way to obtain it, for example, is with $v_1/v_2 \sim 35$.

In conclusion we have shown that it is possible to accommodate the $A_{FB}^t$ 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.

Acknowledgments The work of X.G.H. was supported in part by National Science Council under grants NSC 89-2112-M-002-058, and in part by the Ministry of Education Academic Excellence Project 89-N-FA01-1-4-3. The work of G.V. was supported in part by DOE under contract number DE-FG02-01ER41155.
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