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Violations of universality in a vectorlike extension of the standard model

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Abstract

Violations of universality of couplings in a vectorlike extension of the standard model with three heavy mirror fermion families are considered. The recently observed discrepancies between experiments and the standard model in the hadronic branching fractions R_b and R_c of the Z-boson are explained by the mixing of fermions with their mirror fermion partners.

1 Introduction

The non-perturbative definition of chiral gauge theories in lattice regularization encounters great difficulties [1, 2]. In fact, up to now no acceptable path-integral formulation with exact gauge invariance at finite cut-off is known. The basic obstacle is "fermion doubling", which occurs under very general circumstances [3]. Therefore, it appears natural to look for a formulation based on exact vectorlike gauge invariance [4], which requires the doubling of the light fermion spectrum [5] by mirror fermions [6]. In this vectorlike extension of the standard model the chiral asymmetry of the light fermion spectrum is a low energy phenomenon. At high energies, above the scale of the vacuum expectation value of the Higgs scalar field, the theory becomes symmetric.

A basic feature of the vectorlike extension of the standard model is the mixing of the low-lying fermion states with their heavy mirror fermion partners. The non-zero mixing angles appear in the couplings to the gauge vector bosons and result, among other things, in some small breaking of the universality. In fact, some universality relations are known to be fulfilled to a good accuracy. In order to reproduce them, together with some other basic facts of electroweak phenomenology, one has to choose a particular mixing scheme [5, 7]. This mixing scheme still leaves some room for the breaking of universality at some other places. In particular, as it will be shown in this paper, the discrepancies between experiments and the standard model in the hadronic branching fractions R_b and R_c of the Z-boson can be explained.

Before discussing a simple viable choice of mixings in section 3, first, in the next section, the general properties of fermion mirror fermion mixing will be shortly summarized.

2 The fermion mirror fermion mixing

The mirror partners of light fermions have the same $\mathrm{SU(3)} \otimes \mathrm{SU(2)_L} \otimes \mathrm{U(1)_Y}$ quantum numbers but opposite chiralities. For instance, the right-handed chiral components of mirror leptons form a doublet with $Y = -1$ with respect to $\mathrm{SU(2)_L} \otimes \mathrm{U(1)_Y}$. The mirror partners must be heavy enough in order to forbid their unobserved pair production at present accelerators. Their first observable consequences follow from the mixing with light fermions. The phenomenologically acceptable mixing schemes were discussed in ref. [5] and will be shortly repeated here for the reader's convenience.

In order to discuss the mixing schemes, let us first consider the simplest case of a single fermion (ψ) mirror fermion (χ) pair. The mass matrix on the $(\overline{\psi}_R, \overline{\psi}_L, \overline{\chi}_R, \overline{\chi}_L) \otimes (\psi_L, \psi_R, \chi_L, \chi_R)$ basis is

$$M = \begin{pmatrix} \mu_\psi & 0 & \mu_R & 0 \\ 0 & \mu_\psi & 0 & \mu_L \\ \mu_L & 0 & \mu_\chi & 0 \\ 0 & \mu_R & 0 & \mu_\chi \end{pmatrix} . \qquad (1)$$

Here μ_L, μ_R are the fermion mirror fermion mixing mass parameters, and the diagonal elements are produced by spontaneous symmetry breaking:

$$\mu_\psi = G_{R\psi} v_R \ , \qquad \mu_\chi = G_{R\chi} v_R \ , \tag{2}$$

with the renormalized Yukawa-couplings $G_{R\psi}$, $G_{R\chi}$ and the vacuum expectation value of the Higgs scalar field v_R.

For $\mu_R \neq \mu_L$ the mass matrix M in (1) is not symmetric, hence one has to diagonalize $M^T M$ by $O^T_{(LR)} M^T M O_{(LR)}$, and $M M^T$ by $O^T_{(RL)} M M^T O_{(RL)}$, where

$$O_{(LR)} = \begin{pmatrix} \cos\alpha_L & 0 & \sin\alpha_L & 0 \\ 0 & \cos\alpha_R & 0 & \sin\alpha_R \\ -\sin\alpha_L & 0 & \cos\alpha_L & 0 \\ 0 & -\sin\alpha_R & 0 & \cos\alpha_R \end{pmatrix} \ , \tag{3}$$

and $O_{(RL)}$ is obtained by exchanging the indices $R \leftrightarrow L$. The rotation angles of the left-handed, respectively, right-handed components satisfy

$$\tan(2\alpha_L) = \frac{2(\mu_\chi \mu_L + \mu_\psi \mu_R)}{\mu_\chi^2 + \mu_R^2 - \mu_\psi^2 - \mu_L^2} \ , \qquad \tan(2\alpha_R) = \frac{2(\mu_\chi \mu_R + \mu_\psi \mu_L)}{\mu_\chi^2 + \mu_L^2 - \mu_\psi^2 - \mu_R^2} \ , \tag{4}$$

and the two mass-squared eigenvalues are given by

$$\mu^2_{1,2} = \frac{1}{2} \Big\{ \mu_\chi^2 + \mu_\psi^2 + \mu_L^2 + \mu_R^2$$

$$\mp \Big[(\mu_\chi^2 - \mu_\psi^2)^2 + (\mu_L^2 - \mu_R^2)^2 + 2(\mu_\chi^2 + \mu_\psi^2)(\mu_L^2 + \mu_R^2) + 8\mu_\chi \mu_\psi \mu_L \mu_R \Big]^{\frac{1}{2}} \Big\} \ . \tag{5}$$

The mass matrix itself is diagonalized by

$$O^T_{(RL)} M O_{(LR)} = O^T_{(LR)} M^T O_{(RL)} = \begin{pmatrix} \mu_1 & 0 & 0 & 0 \\ 0 & \mu_1 & 0 & 0 \\ 0 & 0 & \mu_2 & 0 \\ 0 & 0 & 0 & \mu_2 \end{pmatrix} \ . \tag{6}$$

This shows that for $\mu_\psi, \mu_L, \mu_R \ll \mu_\chi$ there is a light state with mass $\mu_1 = \mathcal{O}(\mu_\psi, \mu_L, \mu_R)$ and a heavy state with mass $\mu_2 = \mathcal{O}(\mu_\chi)$. In general, both the light and heavy states are mixtures of the original fermion and mirror fermion. According to (4), for $\mu_L \neq \mu_R$ the fermion-mirror-fermion mixing angle in the left-handed sector is different from the one in the right-handed sector.

In case of three mirror pairs of fermion families the diagonalization of the mass matrix is in principle similar but, of course, more complicated. The strongest constraints on mixing angles between ordinary fermions and mirror fermions arise from the conservation of e-, μ- and τ-lepton numbers and from the absence of flavour changing neutral currents. These constraints can be avoided in a *"monogamous mixing"* scheme, where the structure of the mass matrix is

such that there is a one-to-one correspondence between fermions and mirror fermions. This happens if the family structure of the mass matrix for mirror fermions is closely related to the one for ordinary fermions.

Let us denote doublet indices by $A = 1, 2$, colour indices by $c = 1, 2, 3$ in such a way that the leptons belong to the fourth value of colour $c = 4$, and family indices by $K = 1, 2, 3$. In general, the entries of the mass matrix for three mirror pairs of fermion families are diagonal in isospin and colour, hence they have the form

$$\mu_{(\psi,\chi);A_2c_2K_2,A_1c_1K_1} = \delta_{A_2A_1}\delta_{c_2c_1}\mu^{(A_1c_1)}_{(\psi,\chi);K_2K_1} \ ,$$

$$\mu_{L;A_2c_2K_2,A_1c_1K_1} = \delta_{A_2A_1}\delta_{c_2c_1}\mu^{(c_1)}_{L;K_2K_1} \ , \qquad \mu_{R;A_2c_2K_2,A_1c_1K_1} = \delta_{A_2A_1}\delta_{c_2c_1}\mu^{(A_1c_1)}_{R;K_2K_1} \ . \tag{7}$$

The diagonalization of the mass matrix can be achieved for given indices A and c by two $6 \otimes 6$ unitary matrices $F_L^{(Ac)}$ and $F_R^{(Ac)}$ acting, respectively, on the L-handed and R-handed subspaces:

$$F_L^{(Ac)\dagger}(M^\dagger M)_L F_L^{(Ac)} \ , \qquad F_R^{(Ac)\dagger}(M^\dagger M)_R F_R^{(Ac)} \ . \tag{8}$$

The main assumption of the "monogamous" mixing scheme is that in the family space $\mu_\psi, \mu_\chi, \mu_L, \mu_R$ are hermitian and simultaneously diagonalizable, that is

$$F_L^{(Ac)} = F_R^{(Ac)} = \begin{pmatrix} F^{(Ac)} & 0 \\ 0 & F^{(Ac)} \end{pmatrix} \ , \tag{9}$$

where the block matrix is in (ψ, χ)-space. The Cabibbo-Kobayashi-Maskawa matrix of quarks is given by

$$C^{(c)} \equiv F^{(2c)\dagger}F^{(1c)} \ , \tag{10}$$

independently for $c = 1, 2, 3$. The corresponding matrix with $c = 4$ and $A = 1 \leftrightarrow 2$ describes the mixing of neutrinos, if the Dirac-mass of the neutrinos is non-zero.

A simple example for the "monogamous" mixing is the following:

$$\mu^{(Ac)}_{\chi;K_2K_1} = \lambda^{(Ac)}_\chi \mu^{(Ac)}_{\psi;K_2K_1} + \delta_{K_2K_1}\Delta^{(Ac)} \ ,$$

$$\mu^{(c)}_{L;K_2K_1} = \delta_{K_2K_1}\delta^{(c)}_L \ , \qquad \mu^{(Ac)}_{R;K_2K_1} = \lambda^{(Ac)}_R \mu^{(Ac)}_{\psi;K_2K_1} + \delta_{K_2K_1}\delta^{(Ac)}_R \ , \tag{11}$$

where, as the notation shows, $\lambda^{(Ac)}_\chi$, $\Delta^{(Ac)}$, $\delta^{(c)}_L$, $\lambda^{(Ac)}_R$, $\delta^{(Ac)}_R$ do not depend on the family index.

The full diagonalization of the mass matrix on the $(\psi_L, \psi_R, \chi_L, \chi_R)$ basis of all three family pairs is achieved by the $96 \otimes 96$ matrix

$$\mathcal{O}^{(LR)}_{A'c'K',AcK} = \delta_{A'A}\delta_{c'c}F^{(Ac)}_{K'K} \cdot \begin{pmatrix} \cos\alpha_L^{(AcK)} & 0 & \sin\alpha_L^{(AcK)} & 0 \\ 0 & \cos\alpha_R^{(AcK)} & 0 & \sin\alpha_R^{(AcK)} \\ -\sin\alpha_L^{(AcK)} & 0 & \cos\alpha_L^{(AcK)} & 0 \\ 0 & -\sin\alpha_R^{(AcK)} & 0 & \cos\alpha_R^{(AcK)} \end{pmatrix} \ . \tag{12}$$

$\mathcal{O}^{(RL)}$ is obtained from $\mathcal{O}^{(LR)}$ by $\alpha_L \leftrightarrow \alpha_R$.

In case of $\mu_R = \mu_L$ the left-handed and right-handed mixing angles are the same:

$$\alpha^{(AcK)} \equiv \alpha_L^{(AcK)} = \alpha_R^{(AcK)} \ . \tag{13}$$

In ref. [5] only this special case was considered. The importance of the left-right-asymmetric mixing was pointed out in ref. [7], where the constraints arising from the measured values of anomalous magnetic moments were investigated. It turned out that strong constraints arise for the product $\alpha_L^{(l)}\alpha_R^{(l)}$ $(l = e, \mu)$, but the individual values $\alpha_L^{(l)}$ and $\alpha_R^{(l)}$ are much less restricted. In case of the L-R asymmetric mixing these constraints can be satisfied, if either the left- or right-handed mixing exactly vanishes (or is very small): $\alpha_L^{(l)} \simeq 0$ or $\alpha_R^{(l)} \simeq 0$.

Another important observation in ref. [7] is the relation

$$\mu_L^{(c)} = -\sin\alpha_R^{(AcK)} \cos\alpha_L^{(AcK)} \mu_1^{(AcK)} + \sin\alpha_L^{(AcK)} \cos\alpha_R^{(AcK)} \mu_2^{(AcK)} \ , \tag{14}$$

which implies for small mixing angles and large mass differences between the two mixed fermion states

$$\mu_L^{(c)} \simeq \alpha_L^{(AcK)} \mu_2^{(AcK)} \qquad (A = 1, 2;\ K = 1, 2, 3) \ . \tag{15}$$

Hence the left-handed mixing angles are inversely proportional to the heavy fermion masses. In addition, in case of the simple mass matrix pattern in (11), the left-hand side is independent of the isospin and family index. No such constraints exist for the right handed mixing angles, therefore the choice

$$\mu_L^{(c)},\ \alpha_L^{(AcK)} \simeq 0 \ , \qquad |\alpha_L^{(AcK)}| \ll |\alpha_R^{(AcK)}| \tag{16}$$

leaves more freedom for large mixing than the other possibility, namely, $\mu_R^{(Ac)},\ \alpha_R^{(AcK)} \simeq 0$.

3 Explanation of R_b and R_c

In order to discuss the couplings of the electroweak vector bosons (A, W, Z), let us consider the form of the electroweak currents. The electroweak interaction can be written in general as

$$eJ_Q(x)_\mu A(x)_\mu + g\left[J_L^-(x)_\mu W^+(x)_\mu + J_L^+(x)_\mu W^-(x)_\mu\right] + \frac{e}{\sin\theta_W \cos\theta_W} J_Z(x)_\mu Z(x)_\mu \ . \tag{17}$$

Here θ_W is the Weinberg-angle ($\sin\theta_W \equiv e/g$, $\sin^2\theta_W \simeq 0.225$).

Let us denote the fermion fields corresponding to the mass eigenstates by $\xi^{(AcK)}(x)$ for the light states and $\eta^{(AcK)}(x)$ for the heavy states, respectively. These are linear combinations of the fermion ($\psi^{(AcK)}(x)$) and mirror fermion ($\chi^{(AcK)}(x)$) fields, as discussed in the previous section. The electromagnetic current $J_Q(x)_\mu$ in (17) has the same form in terms of ξ and η as in terms of ψ and χ:

$$J_Q(x)_\mu = \sum_{A,c,K} Q^{(Ac)} \left\{\overline{\xi}^{(AcK)}(x)\gamma_\mu \xi^{(AcK)}(x) + \overline{\eta}^{(AcK)}(x)\gamma_\mu \eta^{(AcK)}(x)\right\} \ , \tag{18}$$

where $Q^{(Ac)}$ denotes the electromagnetic charge. The neutral current $J_Z(x)_\mu$ can be expressed as

$$J_Z(x)_\mu = \sin^2\theta_W J_Q(x)_\mu - J_{L3}(x)_\mu \ , \tag{19}$$

with the third isospin component of the left-handed current $J_{L3}(x)_\mu$. This latter depends on the mixing angles but is still diagonal in isospin:

$$\begin{aligned} J_{L3}(x)_\mu = \frac{1}{2}\sum_{A,c,K}(-1)^{1+A}\Big\{ & \cos^2\alpha_L^{(AcK)}\overline{\xi}_L^{(AcK)}(x)\gamma_\mu\xi_L^{(AcK)}(x) + \sin^2\alpha_R^{(AcK)}\overline{\xi}_R^{(AcK)}(x)\gamma_\mu\xi_R^{(AcK)}(x) \\ & + \sin^2\alpha_L^{(AcK)}\overline{\eta}_L^{(AcK)}(x)\gamma_\mu\eta_L^{(AcK)}(x) + \cos^2\alpha_R^{(AcK)}\overline{\eta}_R^{(AcK)}(x)\gamma_\mu\eta_R^{(AcK)}(x) \\ & + \cos\alpha_L^{(AcK)}\sin\alpha_L^{(AcK)}\left[\overline{\xi}_L^{(AcK)}(x)\gamma_\mu\eta_L^{(AcK)}(x) + \overline{\eta}_L^{(AcK)}(x)\gamma_\mu\xi_L^{(AcK)}(x)\right] \\ & - \cos\alpha_R^{(AcK)}\sin\alpha_R^{(AcK)}\left[\overline{\xi}_R^{(AcK)}(x)\gamma_\mu\eta_R^{(AcK)}(x) + \overline{\eta}_R^{(AcK)}(x)\gamma_\mu\xi_R^{(AcK)}(x)\right]\Big\} \ . \end{aligned} \tag{20}$$

Finally, the charged current contains also the CKM-matrices defined in (10):

$$\begin{aligned} J_L^+(x)_\mu = \frac{1}{\sqrt{2}}\sum_{c,K_1,K_2} C^{(c)}_{K_1K_2}\cdot \\ \cdot\Big\{ & \cos\alpha_L^{(1cK_1)}\cos\alpha_L^{(2cK_2)}\overline{\xi}_L^{(2cK_2)}(x)\gamma_\mu\xi_L^{(1cK_1)}(x) + \sin\alpha_R^{(1cK_1)}\sin\alpha_R^{(2cK_2)}\overline{\xi}_R^{(2cK_2)}(x)\gamma_\mu\xi_R^{(1cK_1)}(x) \\ & + \sin\alpha_L^{(1cK_1)}\sin\alpha_L^{(2cK_2)}\overline{\eta}_L^{(2cK_2)}(x)\gamma_\mu\eta_L^{(1cK_1)}(x) + \cos\alpha_R^{(1cK_1)}\cos\alpha_R^{(2cK_2)}\overline{\eta}_R^{(2cK_2)}(x)\gamma_\mu\eta_R^{(1cK_1)}(x) \\ & + \cos\alpha_L^{(1cK_1)}\sin\alpha_L^{(2cK_2)}\overline{\eta}_L^{(2cK_2)}(x)\gamma_\mu\xi_L^{(1cK_1)}(x) + \sin\alpha_L^{(1cK_1)}\cos\alpha_L^{(2cK_2)}\overline{\xi}_L^{(2cK_2)}(x)\gamma_\mu\eta_L^{(1cK_1)}(x) \\ & - \sin\alpha_R^{(1cK_1)}\cos\alpha_R^{(2cK_2)}\overline{\eta}_R^{(2cK_2)}(x)\gamma_\mu\xi_R^{(1cK_1)}(x) - \cos\alpha_R^{(1cK_1)}\sin\alpha_R^{(2cK_2)}\overline{\xi}_R^{(2cK_2)}(x)\gamma_\mu\eta_R^{(1cK_1)}(x)\Big\} \ . \end{aligned} \tag{21}$$

Note that the CKM-structure in the charged currents is exactly conserved only in the case of

$$\alpha_L^{(AcK_1)} = \alpha_L^{(AcK_2)} \ , \qquad \alpha_R^{(AcK_1)} = \alpha_R^{(AcK_2)} \ , \tag{22}$$

that is for universal mixing angles independent of the family index. In this interesting special case both charged and neutral currents satisfy exact universality with respect to the family index, if quarks and leptons are considered separately.

There is, however, still a possibility to break universality of the couplings between quarks and leptons. Such universality relations are not as strong as for leptons and quarks separately, because there the strong coupling α_s appears due to QCD radiative corrections. For instance, in the Z-boson widths the QCD correction factors for hadronic final states originating from light ($m_q \simeq 0$) quarks have the form $1 + K_{QCD}$ with [8]

$$K_{QCD} = \frac{\alpha_s}{\pi} + 1.409\left(\frac{\alpha_s}{\pi}\right)^2 - 12.77\left(\frac{\alpha_s}{\pi}\right)^3 + \dots \ . \tag{23}$$

Therefore the existing few percent uncertainty in $\alpha_s(M_Z)$ results in an overall uncertainty of the order of percents in the universality relation of the hadronic versus leptonic Z decay widths.

Another observation about the currents in eqs. (20-21) is that on the light states, for small mixing angles, the admixtures of $(V+A)$-type are proportional to α_R^2 (in amplitude). This means that both neutral and charged currents are strongly dominated by the $(V-A)$ couplings, in accordance with observations. Moreover, choosing the ordering in (16), and in addition requiring no (or very small) mixings for one of the members of the isospin doublets, for instance,

$$\alpha_R^{(\nu)}, \alpha_R^{(d)} \simeq 0 \ , \tag{24}$$

the charged currents on the light states become perfectly $(V-A)$.

After these preparations let us now turn in more detail to the question of violations of universality in Z-boson decay. According to eqs. (19-20) the usual coupling parameters g_V^f and g_A^f for the vector and axialvector couplings of the Z-boson become

$$g_V^f = T_{3L}\left[\cos^2\alpha_L^{(f)} + \sin^2\alpha_R^{(f)}\right] - 2Q_f \sin^2\theta_W \ ,$$

$$g_A^f = T_{3L}\left[\cos^2\alpha_L^{(f)} - \sin^2\alpha_R^{(f)}\right] \ , \tag{25}$$

where $T_{3L} = \pm\frac{1}{2}$ is the third component of the left-handed isospin, $\alpha_{L,R}^{(f)} \equiv \alpha_{L,R}^{(AcK)}$ denotes the mixing angles and $Q_f \equiv Q^{(Ac)}$ is the electric charge in units of the positron charge. Up to order $\mathcal{O}(\alpha_{L,R}^2)$, we obtain the following tree-level corrections for the $Z \to ff$ widths:

$$\frac{\Gamma_{Z\to ff}}{\Gamma_{Z\to ff}^{(sm)}} = \frac{\cos^4\alpha_L^{(f)} + \sin^4\alpha_R^{(f)} - 4s_W^2|Q_f|(\cos^2\alpha_L^{(f)} + \sin^2\alpha_R^{(f)}) + 8s_W^4Q_f^2}{1 - 4s_W^2|Q_f| + 8s_W^4Q_f^2}$$

$$= \frac{1 - 4s_W^2|Q_f| + 8s_W^4Q_f^2 - \alpha_L^{(f)2}(2 - 4s_W^2|Q_f|) - \alpha_R^{(f)2}4s_W^2|Q_f|}{1 - 4s_W^2|Q_f| + 8s_W^4Q_f^2} + \mathcal{O}(\alpha_{L,R}^4) \ , \tag{26}$$

where the index (sm) denotes the standard model expressions and $s_W^2 \equiv \sin^2\theta_W$. Similarly, the corrections for the usual parameters in the forward-backward asymmetries

$$A_f \equiv \frac{2g_V^f g_A^f}{(g_V^f)^2 + (g_A^f)^2} \tag{27}$$

have the form

$$\frac{A_f}{A_f^{(sm)}} = \frac{\cos^4\alpha_L^{(f)} - \sin^4\alpha_R^{(f)} - 4s_W^2|Q_f|(\cos^2\alpha_L^{(f)} - \sin^2\alpha_R^{(f)})}{\cos^4\alpha_L^{(f)} + \sin^4\alpha_R^{(f)} - 4s_W^2|Q_f|(\cos^2\alpha_L^{(f)} + \sin^2\alpha_R^{(f)}) + 8s_W^4Q_f^2}$$

$$= 1 + \alpha_L^{(f)2}(2 - 4s_W^2|Q_f|)\left[(1 - 4s_W^2|Q_f| + 8s_W^4Q_f^2)^{-1} - (1 - 4s_W^2|Q_f|)^{-1}\right]$$

$$+ \alpha_R^{(f)2}4s_W^2|Q_f|\left[(1 - 4s_W^2|Q_f| + 8s_W^4Q_f^2)^{-1} + (1 - 4s_W^2|Q_f|)^{-1}\right] + \mathcal{O}(\alpha_{L,R}^4) \ . \tag{28}$$

The comparison of these expressions to the high precision leptonic data [9] shows that the leptonic mixing angles are small: $\alpha_{L,R}^{(e,\mu,\tau,\nu)} = \mathcal{O}(10^{-3}) - \mathcal{O}(10^{-4})$. This is in agreement

with earlier observations, for instance, in ref. [7]. In the hadronic sector let us concentrate on the experimental data [9] $\Gamma_{had} = 1744.8(3.0)$ MeV, $R_b = 0.2219(17)$, $R_c = 0.1543(74)$ and $A_{FB}^{(0,b)}/A_{FB}^{(0,c)} = A_b/A_c = 1.38(15)$. In the absence of a full one-loop calculation, we have to use the tree-level expressions. For the Weinberg-angle we consider the values $s_W^2 = 0.230$ and $s_W^2 = 0.225$. The former is a typical value obtained by the one-loop fits within the minimal standard model [9], whereas the latter is obtained from the tree-level relation $s_W^2 = 1 - M_W^2/M_Z^2$. For the strong coupling at $Q^2 = M_Z^2$ we take the three representative values $\alpha_s = 0.123,\ 0.130,\ 0.116$. Assuming that the observed deviations from universality in R_b and R_c are due to a single dominant mixing angle parameter, I tried fits of the above data in several different settings. Representative examples are:

$$A:\ \alpha_R^{(c)2} \neq 0\ , \qquad B:\ \alpha_L^{(u)2} = \alpha_L^{(d)2} = \alpha_L^{(c)2} = \alpha_L^{(s)2} \neq 0\ , \qquad C:\ \alpha_L^{(c)2} = \alpha_L^{(s)2} \neq 0\ . \tag{29}$$

In case of $s_W^2 = 0.230$, $\alpha_s = 0.123 \pm 0.007$ the obtained χ^2 values were in the range $\chi^2 = 6 - 7$ and the best fits for the mixing angles were compatible with zero. (Possible correlations in the data are neglected here.) This reflects the well known fact that the chosen set of data cannot be well described by the minimal standard model. (In fact, taking into account the one-loop electroweak corrections makes the fit even worse, see ref. [9].) In cases B and C better fits are possible with the tree-level value $s_W^2 = 0.225$ and $\alpha_s = 0.130$: for case B $\alpha_L^{(u)2} = \alpha_L^{(d)2} = \alpha_L^{(c)2} = \alpha_L^{(s)2} = 0.0054(10)$, $\chi^2 = 4.7$, and for case C $\alpha_L^{(c)2} = \alpha_L^{(s)2} = 0.011(2)$, $\chi^2 = 3.5$. In these cases the other values of α_s give slightly worse χ^2. In case A $s_W^2 = 0.225$ leads to substantially worse fits than $s_W^2 = 0.230$, namely $\chi^2 = 9 - 12$. This shows that the best case is:

$$C:\ \alpha_L^{(c)2} = \alpha_L^{(s)2} = 0.011(2)\ , \qquad (\chi^2 = 3.5)\ . \tag{30}$$

The best fit for case B is somewhat worse than for C, but still better than the standard model fit.

4 Discussion

As it has been shown in the previous section, the violations of universality in R_b and R_c at LEP can be explained in a vectorlike extension of the standard model by the mixing of the low-lying fermion states with the heavy mirror fermion partners. A simple choice of the dominant mixing angles for this is given in eq. (30), where the universality violations are mainly due to the equal left-handed mixings of the second family quarks. By this choice the violations of universality at tree level are minimized but, of course, other mixing schemes with several non-negligible mixing angles cannot be excluded at present. (For some other ways to explain R_b and R_c by mixing to heavy states see ref. [10].)

The analysis in the present paper relies on the tree approximation. For a more accurate treatment a full calculation of one-loop radiative corrections is necessary, including also vertex- and box-corrections. First steps in this direction were done in ref. [11], where the one-loop

effects in vector boson propagators were considered and an overall fit of the experimental data was performed, taking into account radiative corrections.

The better description of the hadronic branching ratios of the Z-boson is a hint for the vectorlike extension of the standard model considered here. The final check is, of course, to find the heavy mirror fermions.

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