FLAVOUR-CHANGING Z^0 DECAYS ☆

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We compute neutral Higgs and Z^0-boson couplings to quarks in a model with three quark–lepton families and one mirror family. Mixing between ordinary and mirror quarks leads to flavour-changing couplings which are negligible for the known quarks. However, t-quark couplings are sizeable and the branching ratio for the decay $Z^0 \to (\bar{t}c) + (t\bar{c})$ can be as large as 10^{-4}.

The observed strong suppression of flavour-changing neutral curents led to the hypothesis of "natural flavour conservation" [1] which requires that all quarks of a given charge and chirality have the same weak isospin and obtain their mass from a single source. This implies that flavour-changing couplings of neutral Higgs and gauge bosons are absent at tree level, and that also one-loop corrections are suppressed due to GIM [2] cancellations. So far experimental bounds on flavour-changing couplings only exist for quarks whose masses are much smaller than the Z^0-boson mass. It is conceivable that the t-quark, whose mass [3] is already known to be of the same order of magnitude as the Z^0-boson mass, does have significant flavour-changing couplings. In the case $m_t < m_Z$ this can be tested in the near future by studying decays of Z^0-bosons produced at SLC and LEP.

In the standard model rates for flavour-changing Z^0-decays [4] are too small to be observed. This is also the case for all extensions in which "natural flavour conservation" is maintained, such as models with two Higgs doublets [5] or supersymmetric models [6]. However, in the present note we want to point out that the observed strong suppression of flavour-changing neutral currents could just be due to the smallness of the masses of the quarks involved, and that for very heavy quarks flavour diagonal and non-diagonal couplings could have comparable strength. As an example we will present a model with three quark–lepton families and one mirror family in which the predicted strength of non-diagonal Z^0-boson couplings to quarks with masses m_i and m_j is $O(m_i m_j / m_W^2)$, where m_W denotes the W-mass. As we shall see, these couplings satisfy all existing experimental bounds on flavour-changing neutral currents and lead only for the decay $Z^0 \to \bar{t}c(t\bar{c})$ to a sizeable deviation from standard model predictions.

Mass dependent, flavour non-diagonal couplings of neutral bosons have previously been discussed in the literature, especially by del Aguila and Bowick [7] and by Fishbane, Norton and Rivard [8]. These authors consider a "see-saw" type mechanism where "heavy" quarks, which form a vector representation of the weak gauge group $SU(2)_W \times U(1)_Y$, mix with the "light" ordinary quarks. The scheme presented in this paper differs from this "see-saw" mechanism in two ways: The additional quarks form a chiral representation of

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$SU(2)_W \times U(1)_Y$, and hence their masses are also generated through spontaneous symmetry breaking, i.e., they are expected to be of order m_W, like the t-quark mass.

Since the t-quark is much heavier than the down-quarks of all three families, the largest flavour non-diagonal couplings will occur in the up-quark sector. In the following we will therefore ignore the down-quark mass matrix and hence not address the possible connection [9] between quark masses and Kobayashi–Maskawa mixing angles.

Let us consider as an explicit example a unified model with SO(10) fermion representation $\boldsymbol{F} = 3\times 16 + 16^* + 3\times 10$ plus additional singlets, which occurs in supersymmetric σ-models based on E_8 [10]. The quark content of the three families, the mirror family and the three ten-plets reads:

$$16:\quad q_{Li} = \begin{pmatrix} u \\ d \end{pmatrix}_{Li}, \quad u_{Ri}, \quad d_{Ri}, \quad i = 1, ..., 3,$$

$$16^*:\quad Q_R = \begin{pmatrix} U \\ D \end{pmatrix}_R, \quad U_L, \quad D_L,$$

$$10:\quad \tilde{d}_{L\alpha}, \quad \tilde{d}_{R\alpha}, \quad \alpha = 1, ..., 3, \tag{1}$$

where q_L and Q_R denote SU(2) doublets, and u_R, d_R, U_L, D_L, $\tilde{d}_L$ and $\tilde{d}_R$ are SU(2) singlets. The E_8 σ-models possess global symmetries which in general forbid direct fermion mass terms. Mass generation through spontaneous symmetry breaking requires at least three multiplets of scalar fields: one SU(2) doublet ϕ, and two SU(2) singlets Φ_1 and Φ_2, which are contained in the set of scalar fields of the σ-model. The corresponding, most general Yukawa interactions are given by

$$L_Y = \Gamma^{(u)}_{ij} \bar{q}_{Li} \tilde{\phi} u_{Rj} + \Gamma^{(d)}_{ij} \bar{q}_{Li} \phi d_{Rj} + \Gamma^{(U)} \bar{Q}_R \tilde{\phi} U_L + \Gamma^{(D)} \bar{Q}_R \phi D_L + \tilde{\Gamma}^{(d)}_{\alpha\beta} \bar{\tilde{d}}_{L\alpha} \Phi_1 \tilde{d}_{R\beta} + \gamma^{(Q)}_i \bar{q}_{Li} \Phi_2 Q_R$$
$$+ \gamma^{(U)}_i \bar{u}_{Ri} \Phi_2 U_L + \gamma^{(D)}_i \bar{d}_{Ri} \Phi_2 D_L + \gamma^{(d)}_{i\alpha} \bar{q}_{Li} \phi \tilde{d}_{R\alpha} + \tilde{\gamma}^{(d)}_{\alpha i} \bar{\tilde{d}}_{L\alpha} \Phi_i d_{Ri} + \text{h.c.} \tag{2}$$

This lagrangian is invariant under the following two discrete symmetries which forbid direct mass terms:

$$P_M:\quad Q_R \to -Q_R, \quad U_L \to -U_L, \quad D_L \to -D_L, \quad \Phi_2 \to -\Phi_2, \quad \tilde{P}:\quad \tilde{d}_{L\alpha} \to -\tilde{d}_{L\alpha}, \quad \Phi_1 \to -\Phi_1. \tag{3a,b}$$

P_M and $\tilde{P}$ act trivially on all other fields.

The scalar fields ϕ, Φ_1 and Φ_2 acquire vacuum expectation values which depend on the specific form of the Higgs potential, and which will all be of order m_W in the simplest class of models. This generates mass matrices for up- and down-quarks whose diagonalization yields flavour diagonal and off-diagonal charged currents, given by the Kobayashi–Maskawa (KM) matrix, as well as neutral currents expressed in terms of mass eigenstates. Since the lagrangian (2) contains quarks with the same charge and chirality which have different weak isospin and couple to more than one Higgs field, the criteria of "natural flavour conservation" [1] are not fulfilled and we expect flavour non-diagonal couplings of neutral bosons. However, as we will see, these couplings are "naturally" proportional to the masses of the various fermion pairs and hence very small except for couplings involving the t-quark. In the following we will therefore focus on the up-quark mass matrix and ignore the down-quark sector where flavour non-diagonal couplings are expected to be small.

After spontaneous symmetry breaking one obtains from the lagrangian (2) an up-quark mass matrix of the form

$$L^{(u)}_M = \bar{u}_{Li} M_{ij} u_{Rj} + \bar{u}_{Li} M_{i4} U_R + \bar{U}_L M_{4i} u_{Ri} + \bar{U}_L M_{44} U_R + \text{h.c.} \tag{4}$$

The matrix M_{ij}, written in its diagonal basis, contains the mass hierarchy $m_u \ll m_c \ll m_t$, which is due to an assumed hierarchy of Yukawa couplings. Since experimental bounds on the t-quark mass already imply $m_t = O(m_W)$, and the mass m_T of the mirror up-quark is also generated through $SU(2)_W \times U(1)_Y$ breaking, one has $m_t \sim m_T \sim m_W$. For the mixing terms M_{4i} and M_{i4} the simplest assumption is that they contain the same hierarchy as the matrix M_{ij}. This leads to the following ansatz for the mass matrix:

$$M=\begin{pmatrix} m_1 & 0 & 0 & \mu_1 \\ 0 & m_2 & 0 & \mu_2 \\ 0 & 0 & m_3 & \mu_3 \\ \mu_1 & \mu_2 & \mu_3 & m_4 \end{pmatrix} \tag{5a}$$

with

$$m_1,\, \mu_1 \ll m_2,\, \mu_2 \ll m_3,\, \mu_3,\, m_4\,. \tag{5b}$$

For simplicity, all mass parameters are assumed to be real.

Because of the assumed hierarchy (5b) the matrix (5a) is easily diagonalized. A straightforward calculation yields for the corresponding orthogonal transformation:

$$U=\begin{pmatrix} 1 & -\mu_1(\mu_2/m_2)(c^2M_1+s^2M_2)/M_1M_2 & \mu_1 s/M_1 & \mu_1 c/M_2 \\ \mu_1(\mu_2/m_2)(c^2M_1+s^2M_2)/M_1M_2 & 1 & \mu_2 s/M_1 & \mu_2 c/M_2 \\ -\mu_1 cs(M_2-M_1)/M_1M_2 & -\mu_2 cs(M_2-M_1)/M_1M_2 & c & -s \\ -\mu_1(c^2M_1+s^2M_2)/M_1M_2 & -\mu_2(c^2M_1+s^2M_2)/M_1M_2 & s & c \end{pmatrix} \tag{6a}$$

with

$$c=\cos\Theta,\quad s=\sin\Theta,\ \tan 2\Theta=\frac{2\mu_3}{m_3-m_4},\quad M_{1,2}=\tfrac{1}{2}(m_3+m_4)\pm\sqrt{\mu_3^2+\tfrac{1}{4}(m_3-m_4)^2}\,. \tag{6b}$$

Here we have only kept the leading term for each matrix element.

Given the matrix U, which connects weak and mass eigenstates, the Higgs and Z^0-boson couplings of mass eigenstates are easily obtained. Following Langacker and London [11], who recently analyzed systematically flavour-conserving mixings between ordinary and exotic fermions, we write the mixing matrix in block form:

$$U=\begin{pmatrix} A & E \\ F & G \end{pmatrix}, \tag{7}$$

where A and F act on sequential up-quarks, and E and G act on the mirror up-quark. The Z^0-boson couplings are then given by (cf. ref. [11])

$$J_Z^\mu=\bar{u}_{Li}\gamma^\mu u_{Li}+\bar{U}_R\gamma^\mu U_R-2\sin^2\Theta_w J_{EM}^\mu+\bar{u}_i\gamma^\mu\gamma_5(F^\dagger F)_{ij}u_j+\bar{u}_i\gamma^\mu\gamma_5(F^\dagger G)_i U+\text{h.c.}-\bar{U}\gamma^\mu\gamma_5(E^\dagger E)U\,, \tag{8}$$

where Θ_w and J_{EM}^μ are the weak mixing angle and the electromagnetic current. With $\mu_i=O(m_i)$ and $m=O(M_{1,2})$ one easily obtains from eqs. (6) and (7):

$$(F^\dagger F)_{ij}=O(m_i m_j/m^2),\quad (F^\dagger G)_i=O(m_i/m),\quad E^\dagger E=O(1)\,. \tag{9a,b,c}$$

From (9a) and (9b) we conclude that the couplings of the top-quark and the mirror quark are of comparable strength.

The couplings of the neutral scalars ϕ^0 and Φ_2, are given by

$$K_D=U^T M_D U,\quad K_{ND}=U^T M_{ND} U\,, \tag{10}$$

where M_D and M_{ND} denote the diagonal and off-diagonal parts of the mass matrix M given in eq. (5). From eqs. (6) and (10) one obtains for the Higgs interactions

$$L_H=\bar{U}_{La}[(m_a\delta_{ab}+\lambda_{ab}^{(0)})\phi^0+\lambda_{ab}^{(2)}\Phi_2]U_{Rb}+\text{h.c.},\quad U_a=(u_i,U)\,, \tag{11a}$$

where

$$\lambda_{ab}^{(0)}=O(m_a m_b/m),\quad \lambda_{ab}^{(2)}=O(m_a m_b/m)\,. \tag{11b,c}$$

Table 1
The limits on flavour-changing Z^0-boson couplings defined in eq. (8); $\lambda_{ij}=(F^+F)_{ij}$. The experimental upper limits are taken from ref. [11]. The theoretical estimates are based on eq. (9) with current quark masses taken from ref. [13] and $m_t>45$ GeV.

Process	Coupling	Upper limit	Theory
K^0–$\bar{K}^0$	$\lvert\lambda_{ds}\rvert$	3×10^{-4}	10^{-7}
$K_L\to\mu^+\mu^-$	$\lvert\lambda_{ds}\rvert$	6×10^{-5}	10^{-7}
D^0–$\bar{D}^0$	$\lvert\lambda_{uc}\rvert$	5×10^{-4}	10^{-6}
B_d–$\bar{B}_d$	$\lvert\lambda_{db}\rvert$	4×10^{-4}	10^{-5}
$B_s\leftrightarrow\ell^+\ell^-X$	$\lvert\lambda_{sb}\rvert$	1×10^{-2}	10^{-4}
$Z^0\to t\bar{c}(\bar{t}c)$	$\lvert\lambda_{ct}\rvert$	–	10^{-2}

Eqs. (9) and (11) show that the flavour-changing Higgs and Z^0-boson couplings have the same mass dependence, which follows from the assumed mass matrix (5). Different mass matrices will lead to different flavour-changing neutral currents.

The phenomenologically most interesting implication of the flavour-changing couplings (9) and (11) is the rather large branching ratio for the decay $Z^0\to\bar{t}c(t\bar{c})$. With $m=O(m_W)$ one obtains from eq. (8):

$$\mathrm{BR}(Z^0\to t\bar{c}(\bar{t}c))\sim\frac{8}{1+(1-\frac{8}{3}\sin^{-2}\Theta_w)^2}\frac{m_c^2m_t^2}{m_W^4}(1-m_t^2/m_Z^2)^2(1+m_t^2/2m_Z^2)\times\mathrm{BR}(Z^0\to c\bar{c})\,. \quad (12)$$

Using $\mathrm{BR}(Z^0\to c\bar{c})\simeq0.10$ one obtains $\mathrm{BR}(Z^0\to\bar{t}c(t\bar{c}))\sim10^{-5}$–$10^{-4}$ for $m_t=(45$–$80)$ GeV. With the expected number of 10^6–10^7 Z^0-bosons per year at LEP [12] and the clean signatures of heavy top decays these branching ratios should be detectable.

Flavour-changing processes not involving top-quarks are too rare to be observable within our model. In table 1 upper limits on various flavour-changing couplings are compared with the predictions of eq. (9). Here we have assumed that the flavour non-diagonal down-quark couplings are also given by eq. (9). Although the mixing in the down-quark sector will be more complicated due to the additional quarks $\tilde{d}_{L,R}$ (cf. eq. (1)), this is likely to be the case since the ordinary down-quarks are all much lighter than the new ones. Hence the mixing is essentially of "see-saw" type which is known [7,8] to lead for the light quarks to couplings of the form (9).

The theoretical predictions in table 1 are about two orders of magnitude smaller than the experimental bounds. This means that flavour non-diagonal couplings could even be larger than those given by eq. (9), which would also lead to a branching ratio for $Z^0\to\bar{t}c(t\bar{c})$ larger than the prediction of eq. (12). This would be the case, for instance, in a Fritzsch-type ansatz [14] for the mass matrix which yields flavour-changing Z^0-couplings near the present experimental limits (cf. table 1) and predicts $\lambda_{ct}\sim10^{-1}$.

According to eq. (9) flavour-changing couplings of the t-quark and the mirror T-quark are of the same order of magnitude. Hence, if $m_T<m_Z-m_c$, the decay $Z^0\to\bar{T}c(T\bar{c})$ is kinematically allowed and has also a branching ratio given by eq. (12) with m_t replaced by m_T.

We conclude that the decay $Z^0\to\bar{t}c(t\bar{c})$ provides a rather unique test of the concept of "natural flavour conservation". the observation of this decay mode with a branching ratio larger than 10^{-6} would require flavour-changing couplings at tree-level. In the framework of a renormalizable theory of electroweak interactions and fermion mass generation this would support the existence of new quarks with weak isospin quantum numbers different from those of ordinary quarks, as well as the existence of additional Higgs scalars. The model presented in this paper illustrates how these additional fields can give a sizeable $(t\bar{c})$ coupling of the Z^0-boson without violating the strong experimental bounds on flavour-changing couplings of the Z^0-boson to other quark pairs.

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