ANALYSIS OF NEUTRAL CURRENTS IN SUPERSTRING INSPIRED MODELS

F. del Aguila^{*)}, G.A. Blair^{")}, M. Daniel⁺⁾ and G.G. Ross^{")}

CERN - Geneva

A B S T R A C T

We fit existing neutral current data to (E_6) models inspired by superstrings, concentrating on the case of diagonal neutral currents. Using a convenient parametrization (which includes all popular models as particular cases) we conclude that $\sin^2\theta_W \sim 0.225$ and $M_Z \sim 92.5$ GeV but we cannot exclude in a model independent way new gauge bosons with masses ~ 100 GeV. We compare fits to different sets of data and discuss particular models.

^{*)} On leave from Dept. de Fisica Teorica, UAB, Bellaterra, Barcelona, Spain.

^{")} Department of Theoretical Physics, University of Oxford, OX1 3NP, U.K.

⁺⁾ Rutherford Appleton Laboratory, Chilton, Didcot, Oxon, OX11 0QX, U.K.

1. Introduction

Superstrings^[1] offer the exciting possibility of a finite unified theory of all four fundamental interactions. However, their ultimate test will rely on measurements made in the laboratory at scales far below the Planck scale where the full structure of the superstring is manifest. Extraction of experimental predictions from the superstring is made difficult by the uncertainties in the compactification^[2] needed to go from the ten dimensional world of the underlying superstring to the four dimensional world of experimental physics. After allowing for the uncertainties, the most definite superstring prediction for the low energy world is the gauge group structure. Realistic superstring theories give rise to an effective low energy gauge group which is contained in the exceptional group E_6 ^[3]. After symmetry breaking this leads to the standard model plus, possibly, new low mass gauge bosons contained in E_6 which may give rise to observable effects.

Unfortunately, there is no definite prediction for which, if any, of the new gauge bosons within E_6 will remain light so, to give model independent predictions, it is necessary to parametrise the most general gauge boson structure within E_6 . It is the purpose of this paper to provide such a parametrisation for the neutral bosons and to apply it to an analysis of data sensitive to neutral-weak currents. The data set analysed consists of 84 data points: following ref. [4] 18 points are taken from νq interactions and 7 from νl interactions, 9 from atomic parity violation as given in ref. [5], 11 from the eD asymmetry^[6], 1 from μC asymmetry^{*[7]}, 12 from $\mu\bar{\mu}$, 8 from $\tau\bar{\tau}$, 8 from $c\bar{c}$, 8 from $b\bar{b}$ asymmetries^[5], and 2 from the UA1 and UA2 values for the Z mass^[8]. We do not include production data from the collider. However we do use the UA1 and UA2 values for the W mass, but mainly as a constraint. We will concentrate on the most likely case of abelian neutral interactions, for which the general effective Lagrangian (which includes previously considered models as particular cases) involves 11 free parameters when no constraint on the neutral gauge boson mass

* We use the experimental value b (in ref [7]) for a beam energy of 120 GeV.

matrix is imposed. If, as it is in the case of superstring inspired models, the Higgs sector involves only fundamental representations of E_6 the number of free parameters reduces to 9. Since loop corrections are small our analysis will involve only tree level contributions* (supersymmetric partners then play no role). [Apart from the Z mass we use in our fit the phenomenological parameters defined in ref. [9,10]. We give below their expressions for the models we consider].

In Section 2 we discuss the general parametrization of extended electroweak models in E_6 , concentrating on the case of diagonal neutral currents. We count the parameters needed to fix the mixing and masses of the gauge bosons. We also discuss the angles needed to assign the known fermions in the fundamental representation of E_6 , although they will not be included in the fits we present in Section 3, where we take the usual assignments for the known fermions corresponding to zero mixing. This is what it has been assumed in all popular models [11,12,13] discussed previously.

2. Parametrization of E_6 models

In E_6 [3] with the standard embedding of $SU_3^C \times SU_2^L \times U_1^Y$ [14] there are 8 $SU_3^C \times U_1^{EM}$ neutral gauge bosons. The adjoint representation of E_6 decomposes under $SU_3^C \times SU_3^L \times SU_3^R$ according to

$$78 = (8,1,1) + (1,8,1) + (1,1,8) + (3,\bar{3},\bar{3}) + (\bar{3},3,3). \quad (2.1)$$

The neutral generators which are colour singlets are contained in the set of generators of SU_3^L and SU_3^R , $(1,8,1)$ and $(1,1,8)$ respectively. To each set correspond four charged gauge bosons and four neutral ones. Two of the neutral bosons are associated with the diagonal generators of $SU(3)$. The other two neutral generators are off-diagonal generators of a $SU(2)$ neutral subgroup of $SU(3)$. So, in general, neutral currents can be mediated by 8 gauge bosons. One of these is the photon and does not mix, but the other

* The important one loop correction to α_{EM} , however, will be taken into account. When needed α_{EM} will be evolved to M_W scale.

seven may have arbitrary mixing and have arbitrary masses (if no assumption on the gauge boson mass generation is made). One, probably the lightest, must have a mass ~ 92 GeV and coupling roughly as expected in the standard model for the Z boson^[8]. As we will discuss later in superstring inspired models with Higgses in fundamental representations of E_6 , the mass matrix for neutral gauge bosons can depend at most on 10 real parameters. In the diagonal case of interest here only five of those parameters are relevant.

In realistic models each fermion generation is assigned to a 27 dimensional fundamental representation of E_6 (we do not discuss possible light $\overline{27}$ components as we are only interested in neutral currents of observed fermions which we assume to lie in 27's). The 27 decomposes under the standard model as follows^{*}

$$\begin{aligned}
 & (u, d_1)_L + d_{2L} + u_L^c + d_{1,2L}^c + (v_{1,2}, e_{1,2})_L + \\
 & (3, 2; \frac{1}{6}) + (3, 1; -\frac{1}{3}) + (\bar{3}, 1; -\frac{2}{3}) + 2(\bar{3}, 1; \frac{1}{3}) + 2(1, 2; -\frac{1}{2}) + \\
 & (e_2^c, \nu_3)_L + e_{1L}^c + \nu_{4,5L} \\
 & (1, 2; \frac{1}{2}) + (1, 1; 1) + 2(1, 1; 0) . \quad (2.2)
 \end{aligned}$$

We have named the different components according to the standard $SU_3^C \times SU_2^L \times U_1^Y$ quantum number assignments of known fermions. As can be observed there is one up quark but two down quarks, two charged leptons, and five neutrinos per family. Under $SU_2^L \times U_1^Y$ the new down quark, charged lepton, and the extra neutrinos have vectorial couplings. In the standard model these new fermions being vector-like are expected to be heavy^[15]. However an extra low energy gauge symmetry can prevent these masses from becoming large. The new fermions in a 27 representation must be heavier than the present threshold of production $\gtrsim 20$ GeV (as we point out later they could already have been produced at the collider)[‡].

^{*} We normalise U_1^Y in the usual way.

[‡] We thank L. Maiani for calling our attention to bounds from W decays.

The general low-energy Lagrangian has two sources of unknown parameters, those coming from the mixing and masses of the residual light gauge bosons and those coming from the mixing and masses of the light fermions. To discuss the former let us start at the Planck scale where all gauge bosons couple with the same coupling. Our analysis applies to the case of (up to four) additional abelian (diagonal) neutral currents.

$$L = g \sum_{i=1}^n J^{\mu i} X_{\mu}^i \quad ; \quad J^{\mu i} = \bar{\psi} T^i \gamma^{\mu} \psi \quad (2.3)$$

where n is the number of neutral gauge bosons and we can choose an arbitrary E_6 basis for the T^i . We specify below our choice. Once E_6 is broken the factors of the unbroken group will renormalise independently due to their different β functions^{*}; so we should rotate to a basis Y_{μ}^i which is multiplicatively renormalised

$$X^i = O^{ij} Y^j \quad (2.4)$$

where O is a $n \times n$ orthogonal matrix. Inserting eq.(2.4) in eq.(2.3) we can now continue to low energies.

$$L = \sum_{i,j} J^{\mu i} O^{ij} g_j Y_{\mu}^j \quad (2.5)$$

* Different U_1 factors may have different β functions after the initial stage of symmetry breaking (flux-breaking) at the compactification scale because the residual light matter fields may be in incomplete E_6 multiplets of the 27 and $\overline{27}$ and thus couple to U_1 factors of E_6 differently.

Here the couplings g_j may differ due to the possibly different renormalisation effects[#]. Since the low-energy group contains SU_2^L , one of the Y^j can always be chosen to be the one associated with T_3^L . This fixes O^{11} . From (2.5) we obtain

$$T_3^L = T^i O^{i1} \quad (2.6)$$

and the significant parameters in O are in a $(n-1) \times (n-1)$ orthogonal submatrix.

Finally we must diagonalise the mass matrix of the Y_μ^j gauge bosons. In terms of the mass eigenstates W_μ^k

$$Y^j = O'^{jk} W^k \quad (2.7)$$

Once again, although O' is a $n \times n$ orthogonal matrix, only a $(n-1) \times (n-1)$ submatrix is needed to describe the most general diagonalisation since we may identify uniquely the photon as a zero mass eigenstate W^1 . That fixes

$$O'^{j1} = g_j^{-1} O^{Tj1} \sqrt{\frac{2}{3}} e \quad (2.8)$$

where e is the positron charge at M_W . Inserting eq.(2.7) in eq.(2.5) and computing the Born graphs gives the effective Lagrangian

$$L^{\text{eff}} = J^\mu O g O' [q^2 - M_D^2]^{-1} O'^T g O^T J_\mu^T \quad (2.9)$$

where M_D^2 is the diagonal mass matrix.

[#]Our general parametrization also applies to the case discussed by Choi and Kim [16] in which U_1 factors after flux breaking have different couplings at the compactification scale. In this case the g_j may differ due both to the different wave function renormalization of the U_1 factors and to their initial values.

However eq. (2.9) is more general than may be expected for E_6 models in which all matter multiplets are in 27 representations of E_6 . This is because a 27 of scalars has only five neutral states which may acquire vacuum expectation values (vevs), breaking the low energy group to $SU_3^C \times U_1^{EM}$. For diagonal neutral currents the mass matrix only involves the modulus of these five vevs i.e. it depends on only five parameters and not the six of a general mass matrix describing three massive vector bosons[‡]. The general mass matrix for the Y_μ^j fields has the form

$$M_{ij}^2 = 2 \sum_k (T \langle \phi^k \rangle_0 g)_i^+ (T \langle \phi^k \rangle_0 g)_j \quad (2.10)$$

where $\langle \phi^k \rangle$ is the vev of a scalar field ϕ^k . Inserting this in eq (2.9) gives

$$\begin{aligned} L^{eff} &= J^\mu [0 \ g^{-1} \ 0' (q^2 - M_D^2) \ 0'^T g^{-1} 0^T]^{-1} J_\mu^T \\ &= J^\mu [0 \ g^{-2} 0^T q^2 - 0 g^{-1} M^2 g^{-1} 0^T]^{-1} J_\mu^T \\ &= J^\mu [0 \ g^{-2} 0^T q^2 - 2 \sum_k (T \langle \phi^k \rangle)^+ (T \langle \phi^k \rangle)]^{-1} J_\mu^T \end{aligned} \quad (2.11)$$

The basis for the generators T may be chosen at will and for convenience we choose

$$\begin{aligned} \tilde{Q} &= \frac{1}{2} \sqrt{\frac{3}{2}} \{ T_3^L + T_3^R + \frac{1}{\sqrt{3}} (Y^L + Y^R) \} \equiv \frac{1}{2} \sqrt{\frac{3}{2}} Q \\ A &= \frac{1}{2} \sqrt{\frac{3}{2}} \{ \frac{1}{\sqrt{3}} (T_3^L + T_3^R) - (Y^L + Y^R) \} \\ B &= \frac{1}{2} \sqrt{\frac{3}{2}} \{ (T_3^L - T_3^R) + \frac{1}{\sqrt{3}} (Y^L - Y^R) \} \\ C &= \frac{1}{2} \sqrt{\frac{3}{2}} \{ \frac{1}{\sqrt{3}} (T_3^L - T_3^R) - (Y^L - Y^R) \} \end{aligned} \quad (2.12)$$

where Q is the charge operator and the action of $T_3^{L,R}$ and $Y^{L,R}$ on the 27 is given in Table 1.

In this basis we find for a typical ϕ^k

[‡] Besides a combination of those five parameters is fixed by the W mass.

$$x_{ij} \equiv 2(T^1 < \phi^k >)^+ (T^1 < \phi^k >)$$

$$= \begin{bmatrix} 0 & 0 & 0 & 0 \\ 0 & (|a_1|^2 + |a_4|^2) & \frac{1}{\sqrt{3}} (|a_1|^2 - |a_4|^2) & 0 \\ 0 & \frac{1}{\sqrt{3}} (|a_1|^2 - |a_4|^2), \frac{1}{3} (4|a_3|^2 + |a_1|^2 + |a_2|^2 + |a_4|^2 + |a_5|^2), \frac{1}{\sqrt{3}} (|a_2|^2 - |a_5|^2) \\ 0 & 0 & \frac{1}{\sqrt{3}} (|a_2|^2 - |a_5|^2) & (|a_2|^2 + |a_5|^2) \end{bmatrix} \quad (2.13)$$

where a_k denotes a vev in a neutral scalar component of a 27 (or $\overline{27}$) corresponding to the positions occupied by the neutral states v_k , $k=1, \dots, 5$ in eq. (2.2) and Table 1. Eq. (2.13) depends only on the five independent moduli of vevs. Adding further multiplets of 27 or $\overline{27}$ will not fill in the zeros of eq. (2.13) so it has the most general possible form. One combination of the vevs will be determined by the charged weak currents because we have

$$M_W^2 = g_1^2 \frac{1}{2} (|a_1|^2 + |a_2|^2 + |a_3|^2) \quad (2.14)$$

where g_1 is the weak coupling constant associated with T_3^L . The matrix X now depends on four free parameters.

The final form of our effective low-energy Lagrangian is

$$L^{\text{eff}} = J^\mu [0 \ g^{-2} \ 0^T \ q^2 - X]^{-1} J_\mu^T \quad (2.15)$$

This form shows that determining the q^2 dependence of the effective Lagrangian is in general essential, if all the parameters are to be determined. If the q^2 dependence were ignored, the 9 parameters (excluding the electric charge) which could be in principle measured would fix in the fully constrained case nine functions of the 9 free parameters: 3 from 0, 2 from g and 4 from X.

The q^2 coefficient $0g^{-2}0^T$ in eq.(2.15) depends in general on 5 free parameters. This follows since 0^{11} is fixed by eq.(2.6) which, with our conventions, reads

$$T_3^L = \frac{1}{2\sqrt{2}} (\sqrt{3}\tilde{Q} + A + \sqrt{3}B + C) \quad (2.16)$$

leaving 3 parameters to fix 0 (see next section for an explicit treatment). Two of the four g 's are also fixed to reproduce the correct electric charge value at M_W and the charged current strength

$$e^2 = \frac{3}{8} (0_{1j}g_j^{-2} 0_{j1}^T)^{-1} \quad (2.17)$$

$$\frac{G_F}{\sqrt{2}} = \frac{g_1^2}{8M_W^2} \quad (2.18)$$

The full set of 5 parameters in the matrix $0g^{-2}0^T$ in general will only be determined, if neutral current measurements at high $q^2 = O(M_Z^2)$ are given.

The effective Lagrangian in eq.(2.15) can now be used to fit available neutral current data. These data are described by standard phenomenological parameters^[9,10]. Using eq.(2.12) and Table 1 they can be written for the effective Lagrangian in eq.(2.15) as a function of

$$B_{ij} = [0g^{-2}0^T q^2 - X]_{ij}^{-1} . \quad (2.19)$$

They are given in Table 2 assuming that the standard fermions correspond to the states as defined in Table 1.

The parametrization of eqs.(2.15) and (2.19) describes the most general neutral current structure following from the string-inspired E_6 gauge structure and subsequent symmetry breaking to the standard model plus abelian factors. Thus it can describe the neutral currents resulting from the exchange of up to three massive neutral vector gauge bosons. The number of parameters needed for this most general case ($n=4$) is 9 namely, 3 describing 0 , 2 describing g , and 4 describing X . In section 3, we give the results of a fit to the available neutral current data using this parameterization. However eqs.(2.15) and (2.19) may not be the most convenient parametrization for phenomenology. Although it clearly describes the independent parameters and shows that their determination requires information at high q^2 , the physical meaning of the parameters is somewhat

remote. In particular the mass matrix following from eq (2.10) has the form

$$M^2 = g \begin{pmatrix} 0 & 0 \\ 0 & 0 \end{pmatrix}^T X \begin{pmatrix} 0 & 0 \\ 0 & 0 \end{pmatrix} g \quad (2.20)$$

To determine the masses of the new neutral gauge bosons requires knowledge of all 9 parameters and diagonalisation of the mass matrix in eq (2.20). A more immediately 'physical' parametrization is given by eqs (2.9) in terms of physical masses. In this case however for $n=4$ there are 11 parameters, namely 3 for 0 , 3 for $0'$, 2 for g , and 3 for M_D^2 . This is two more than needed in eq. (2.15) because it does not contain the constraints following from 27 or $\overline{27}$ breaking of the low-energy group.

This completes our discussion of the general parametrization of the neutral boson structure to be expected in superstring inspired models which lead to additional abelian gauge bosons. The utility of this parametrization is that it allows a model independent fit to the data which, in principle, can determine the pattern of symmetry breaking from the underlying E_6 gauge structure. We will illustrate this in the next section.

Before doing this, however, we must complete our parametrization by including the mixing angles in the fermion sector. As shown in eq (2.2) several states have the same $SU_3^C \times SU_2^L \times U_1^Y$ quantum numbers in a 27 representation, but only one combination in each case corresponds to the observed fermions. This requires new mixing angles. The e_L is fixed by the mixing angle θ_L^e to an $e_{1,2L}$ combination. The state e_L^+ can belong to e_1^C or e_2^C and so we must define a mixing angle θ_R^e when assigning e_L^+ to these states. Finally, to complete the lepton multiplet, the neutrino ν_{eL} must be assigned to $(\nu_1, \nu_2, \nu_3, \nu_4, \nu_5)_L$. The 5×5 orthogonal matrix* describing the general mixing of the five ν states has just four real parameters relevant to low energy neutral current processes describing the ν_{eL} components in ν_i . Thus, six parameters are needed for the multiplet containing the electron although the stringent bounds on electron number conservation will reduce this to five since a significant θ_R^e will require no ν_{eL} component in ν_3 [17].

* We are not concerned here with possible CP violating phases.

We turn now to the quark assignments. The mass eigenstate d_L may be in a mixture of the $d_{1,2L}$ current eigenstates so a general assignment requires a mixing angle ϕ_L^d . To avoid strangeness changing neutral currents due to W exchange the s_L and b_L states - before intergenerational mixing - must be assigned to the same mixture of $s_{1,2L}$ and $b_{1,2L}$, respectively. Quark-lepton universality for the charged weak current requires that the Wud coupling be the same as the W_{ν_e} coupling so a combination of the lepton mixing angles must be constrained to be equal to ϕ_L^d giving no net increase in the number of parameters. Finally a new angle ϕ_R^d must fix the $d_{1,2L}^c$ combination defining d_L^c .

To summarise, the first generation requires six parameters, $\theta_{L,R}^e$, ϕ_R^d and 3 angles for fixing v_L . To this must be added similar angles for the other generations plus intergenerational mixing. However, the observed smallness of flavour changing neutral currents requires that these angles be the same for quarks and charged leptons for all generations, at least to the accuracy relevant when analysing flavour conserving neutral currents. Assuming the same applies to neutral leptons we are left with just the six angles of the first family. They reduce to three if $\phi_L^d = 0$ for, in this case, quark-lepton universality requires v_{eL} be assigned uniquely to $v_{1,2}$ with a mixing angle θ_L^e . We note that the condition $\phi_L^d = 0$ follows in many models which have a different scale for the $(d_{2L}, d_{1,2L}^c)$ [$SU(2)_L$ invariant] mass from the $(d_{1L}, d_{1,2L}^c)$ [$SU(2)_L$ doublet] mass. The same argument leads to $\theta_R^e = 0$, leaving us with just two angles θ_L^e and ϕ_R^d . In our phenomenological analysis we do not include the full set of fermion mixing angles, because the data is insufficient to determine them all. In the reduced case with just θ_L^e and ϕ_R^d these angles enter only in the currents coupled to new gauge bosons and, as the fits prefer these to be heavy, these angles are very poorly determined. Thus for the fits presented here we fix all mixing angles to be zero. If the reader wishes to take fermion mixing into account Table 2 should be modified to include these angles.

3. Phenomenological analysis of low energy neutral current data

In this section we apply our parametrization derived in section 2 to an analysis of the available neutral current data. In addition we constrain the parameters by the observed W boson mass via eq.(2.14), and fix g_1 and α_{EM} to their values at M_W (eqs.(2.17) and (2.18)). We will work only with the parametrization corresponding to zero fermion mixing, because the data are insufficient to determine the general fermion mixing case.

Before discussing our results let us explain how we use eqs.(2.9) and (2.20) together with (2.13,14,16-18) and (2.8) in our fits.

O_{ij} is a 4x4 orthogonal matrix but, cf the discussion leading to eq.(2.6), involves only three independent angles. Thus we write O in the form

$$O_{ij} = \begin{bmatrix} c'_1 c'_2 c'_3 & s'_1 & -s'_2 c'_1 & c'_1 c'_2 s'_3 \\ -s'_1 c'_2 c'_3 & c'_1 & s'_1 s'_2 & -s'_1 c'_2 s'_3 \\ s'_2 c'_3 & 0 & c'_2 & s'_2 s'_3 \\ -s'_3 & 0 & 0 & c'_3 \end{bmatrix} \begin{bmatrix} 1 & 0 \\ 0 & R \end{bmatrix} \quad (3.1)$$

where c'_1 and s'_1 stand for $\cos\theta'_1$ and $\sin\theta'_1$ respectively and R is a general 3x3 orthogonal matrix. Eqs (2.6) and (2.16) fix O_{11} and then θ'_1 , giving

$$O_{ij} = \begin{bmatrix} \frac{\sqrt{3}}{2\sqrt{2}} & -\frac{1}{2} & -\frac{3}{2\sqrt{3}} & -\frac{\sqrt{3}}{2\sqrt{14}} \\ \frac{1}{2\sqrt{2}} & \frac{\sqrt{3}}{2} & -\frac{\sqrt{3}}{2\sqrt{7}} & -\frac{1}{2\sqrt{14}} \\ \frac{\sqrt{3}}{2\sqrt{2}} & 0 & \frac{2}{\sqrt{7}} & -\frac{\sqrt{3}}{2\sqrt{14}} \\ \frac{1}{2\sqrt{2}} & 0 & 0 & \frac{\sqrt{7}}{2\sqrt{2}} \end{bmatrix} \begin{bmatrix} 1 & 0 \\ 0 & R \end{bmatrix} \quad (3.2)$$

with

$$R = \begin{bmatrix} c_3 c_2 & s_1 s_2 c_3 + c_1 s_3 & -c_3 s_2 c_1 + s_1 s_3 \\ -s_3 c_2 & -s_1 s_2 s_3 + c_3 c_1 & s_2 c_1 s_3 + s_1 c_3 \\ s_2 & -c_2 s_1 & c_1 c_2 \end{bmatrix} \quad (3.3)$$

The matrix O given by (3.2) and (3.3), involves 3 free parameters, θ_i . We fix g_j by choosing $g_{3,4}$ as free parameters, then (see (2.18) and (2.17))

$$g_1^2 = \frac{G_F}{\sqrt{2}} 8 M_W^2 \quad (3.4)$$

and

$$g_2^2 = O_{12}^2 \left(\frac{3}{8} \frac{1}{e^2} - O_{11}^2 g_1^{-2} - O_{13}^2 g_3^{-2} - O_{14}^2 g_4^{-2} \right)^{-1}. \quad (3.5)$$

Finally we use X in eq. (2.13) and eq. (2.20) to find O' and M_D^2 as functions of X , O and g . We again constrain X by eq. (2.14), leaving only 4 free parameters. M_W is essentially taken into account as a constraint ((2.18), (2.14)) but for completeness we also do fits with M_W as extra data^[8] and g_1 as free parameter, then we must use

$$|a_1|^2 + |a_2|^2 + |a_3|^2 = (2\sqrt{2} G_F)^{-1}. \quad (3.6)$$

as a constraint on X .

With O , g , O' and M_D^2 we can then use (2.19) to obtain B and Table 2 then parametrises the data. In obtaining B the photon has to be projected out. This is done by summing only over the massive propagators. From eqs (2.9) and (2.19) we obtain

$$B_{ij} = \sum_{k,l=1}^4 \sum_{\alpha=2}^4 O_{ik} g_k O'_{k\alpha} \frac{1}{q^2 - M_{D\alpha}^2} O_{\alpha l}^T g_l O_{lj}^T. \quad (3.7)$$

We are now ready to do the fit of low energy neutral current data to (diagonal) neutral interactions from E_6 . We have done a χ^2 fit of eq. (2.11) to the 84 data points described in the introduction, using the minimization program MINUIT available at CERN^[18]. Except for the last two points corresponding to the Z mass M_1 , the remaining points fix the phenomenological Lagrangian describing small q^2 electroweak interactions. The expressions of the coefficients of the corresponding four fermion operators^[9,10] are given in Table 2 for the general cases we consider. We work at tree level and keep only linear terms in the weak propagators, following standard approximations^[5].

It is worth noting that among the phenomenological coefficients in Table 2 there are linear relations^[10]. In models with SU_5 unification at some scale four of such relations must be fulfilled^[19,20]. In E_6 , in general with the observed fermions as combinations of entries in different SU_5 representations and with extra neutral currents, these relations do not hold. However, for zero fermion mixing (see section 2) three relations exist

$$\begin{aligned}\alpha + 3\gamma + 2g_V &= 0 \\ -\beta + \delta - 2g_A &= 0 \\ \tilde{\alpha} - \tilde{\beta} + 3\tilde{\gamma} + \tilde{\delta} &= 0\end{aligned}\tag{3.8}$$

With our general formulation we can perform a model independent analysis. However for comparison with previous analyses we also do fits to particular cases, which are obtained by fixing some of the new parameters (see below). Our results for the general case are given in Table 3 and for particular cases in Tables 4-8. Thus in Table 4 we fit the standard model for reference. In Table 5 we fit the popular superstring inspired model with an extra Z with equal couplings to fermions carrying the same $SU_3^C \times SU_2^L \times U_1^Y$ quantum numbers^[11,20]. (It is obtained from E_6 when a non-abelian flux symmetry breaking is assumed^[2,21]). In Table 6 we analyse a model^[12] with an extra gauge boson resulting from a large vacuum expectation value along the $\langle v_R \rangle$ entry in the 16 of SO_{10} . Table 7 gives the standard left-right (LR) model fit^[22,23]. Table 8 gives the general fit to models with an intermediate mass scale and usual fermions in the gauge protected sector^[21].

In all the cases we do several fits, using the W mass as a constraint with the values quoted by UA1 and UA2^[8] respectively (hence $\sin^2\theta_W$ is fixed by eq.(3.4)), as well as using them as two extra data points (with $\sin^2\theta_W$ free) (hence the different headings on the Tables).

The analysis constrains the neutral gauge boson mass matrix as discussed in Section 2. However we find the error analysis for the new neutral gauge boson masses is unstable since these masses are complicated functions of our parameters as in eq.(2.20) and the errors are large. For this reason we also do the "physical" fit in terms of unconstrained boson masses which, although more general than the E_6 model allows, does give an upper bound to the allowed mass range for the new gauge bosons.

Finally, we present results for two sets of data points, the complete set of 84 data points and for comparison with previous analyses a subset of 51 which was used in ref [24]. The latter exclude all atomic physics points except one^[25], the μC asymmetry point and the $\tau\tau$, $c\bar{c}$ and $b\bar{b}$ asymmetries. Looking at the χ^2 of both sets in the Tables we see that the presence of new parameters (new physics) does not improve the fits at all when compared with the standard model in Table 4, all having $\chi^2/\text{degree of freedom}$ less than one. This statement is more dramatic when one compares each individual data point, all agree with the standard model predictions. To summarize, the experimental numbers scatter nicely around the standard model predictions for $\sin^2\theta_W \sim .225$ and $M_Z \sim 92.5$ GeV. The χ^2 for the subset of 51 points does not change at all for particular models (Tables 5,6,7). Using those fits and allowing only for one standard deviation we conclude that a new gauge boson ~ 150 GeV cannot be excluded in a model independent way. A word about errors is now appropriate. Because of the non-linearity of the problem the parabolic errors quoted in Tables 5,6,7 are not good. The correct non-parabolic lower bound to the error calculated by MINOS in MINUIT (to one standard deviation) is shown underlined in the above Tables. Now at two standard deviations similar fits of the subset of 51 points allow for new gauge bosons ~ 100 GeV. Tables 3 and 8 show that in the general case even at one standard deviation new gauge bosons with masses ~ 100 GeV cannot be excluded. However, in this case the minimum is in a region with many flat directions and parabolic errors are not very meaningful. Besides, in E_6 , there is a combination of currents corresponding to

$$\frac{1}{2\sqrt{2}} (-2Y^L + \sqrt{3} T_3^R + Y^R) \quad (3.9)$$

which does not couple to the ordinary leptons (when zero fermion mixing is assumed). Then the mass of the gauge boson associated with it cannot be constrained by available neutral current data. In fact in the fit there is a tendency of the extra gauge bosons to align along (3.9).

Although not required by the full data set, it is interesting to ask whether additional neutral currents can give a larger $\mu\bar{\mu}$ asymmetry than the standard model without spoiling the other predictions. To investigate this we have done fits to a subset of 45 data points including only the $\mu\bar{\mu}$ asymmetry points with relatively large value ($Q=34.4, 34.5, 34.6, 34.7, 41.1, 41.5$ GeV)^[5]. In this case the χ^2 for the standard model is 30. The particular cases which have been studied do not improve the fit. However the general case ($\chi^2 = 27$) could explain a large $\mu\bar{\mu}$ asymmetry. This is because of the existence in E_6 of an essentially axial current which, when coupled to a gauge boson ~ 380 GeV, will contribute in increasing the asymmetry. But in this case a third gauge boson ~ 260 GeV is required to cancel the new contributions to the other quantities^{*}. We must stress, however, that presently available data do not require those contributions.

To close the discussion of the data points we want to comment on the possibility that certain combinations of the phenomenological parameters $\alpha, \beta, \gamma, \delta$ in Table 2 could favour departures from the standard model^[20]. However, more recent analyses^[26] give values for α, β, γ and δ (and their combinations) in good agreement with the standard model. Our fits also prefer the standard model.

Let us now discuss the Tables in more detail. They give the values of the free parameters for the best fit (errors correspond to one standard deviation)^{**}. Some of the parameters are very poorly determined by the data and for them (denoted by an asterisk) we use the value predicted by E_6 , assuming no intermediate breaking. When the Higgs sector is fixed we also give the corresponding gauge boson masses and mixings (without errors).

To compare with previous analyses and to understand our limits we have to keep in mind that the errors are calculated by changing χ^2 by one. If a local minimum is reached with a χ^2 better by more than one unity some errors

* We thank P. Langacker for discussions on this possibility.

** When the masses of new gauge bosons go to infinity their effects disappear. In that sense the errors must be understood as giving a lower bound.

can be ridiculously small, showing only the dependence on those parameters of that particular minimum. This is clear in the general case of Table 3, a smaller χ^2 is obtained by requiring new gauge bosons. Because all fits having $\chi^2/\text{degrees of freedom}$ less than one are acceptable, we do not interpret these results as evidence for new physics. The model of Table 5 for the subset of 51 data points has been recently analysed^[24]. Our χ^2 agrees nicely although the errors, which are otherwise compatible, seem in particular to be smaller for the gauge boson masses in our case. This is surprising because as discussed above these errors are obtained using the general parametrization, which gives us upper bound to the errors. This may be due to the difficulty of determining the errors of masses when they are not the free parameters as in the case of ref [24].

4. Discussion

The summary of our analysis is that present data agree with the standard model predictions and by enlarging the model only bounds on new contributions are obtained. In this sense the general fit in Table 3 with the lowest χ^2 is simply a curiosity^{*}. However this fit does illustrate the usefulness of the method, for it allows us to compare all possible models at once and from the fit we could reconstruct the pattern of symmetry breaking from the original E_6 model via eqs (2.3-2.9).

Let us make some final remarks. We have presented a general parametrization of (diagonal) neutral currents in (E_6) superstring inspired models. The results of our fits seem to indicate that a new gauge boson of 0(100) GeV cannot be in general excluded, although heavier masses seem to be preferred (see Tables). One can study the production of these new gauge bosons using their couplings to fermions given by eqs (2.5),(2.7) and Table 1. Perhaps it is more interesting to investigate possible signatures of the fermions (and scalars if supersymmetry protects them) which are chiral only under these new interactions^[15,27], and which will be probably lighter than the corresponding gauge bosons. Much could be learned if one of them were discovered. One cannot exclude the possibility that some of them could have perhaps been already produced at the collider^{**}. In the

^{*} Many of the new parameters prefer the corresponding fixed value in the standard model.

^{**} We thank discussions with L. DiLella on this possibility.

27's of E_6 there are with the standard (e, ν) , new vector-like (E, N) leptons. These should be heavier than ~ 20 GeV (their production threshold at e^+e^-), contributing less than the (τ, ν) to the $(e\nu)$ sample of W decays. They could then explain the 20% disagreement (compatible within errors) of the observed and the predicted product of cross section by branching ratio. Besides being heavier they could better explain some of the observed monojets, a more careful discussion will be presented elsewhere.

Finally one of the motivations of this analysis was to calculate the allowed range of values for $\sin^2\theta_W$ in superstring inspired models for it is a strong constraint for model building^[28]. Our fits confirm that those models must predict $\sin^2\theta_W \sim .225 \pm .015$ (consistent with the results from W production).

Acknowledgements

We acknowledge discussions with B. Adeva, A. De Rujula, L. DiLella, L. Ibanez, P. Langacker, L. Maiani, P. Nilles and F. Zwirner and help on the computing from J. Hernandez and F. Saldana.

Appendix

In this Appendix we give the details of the analysis of the particular models discussed in the text. They can be obtained by specifying the different parameters of the general case and projecting out the extra currents.

(i) In the standard model the two first columns of O_{ij} in (3.1) and (3.2) are fixed

$$O_{ij} = \begin{pmatrix} \frac{\sqrt{3}}{2\sqrt{2}} & \frac{\sqrt{5}}{2\sqrt{2}} & \cdot & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & \cdot & \cdot \\ \frac{\sqrt{3}}{2\sqrt{2}} & -\frac{3}{2\sqrt{10}} & \cdot & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & \cdot & \cdot \end{pmatrix} \quad (A.1)$$

The extra ones are irrelevant and any summation should only include the first two columns. In this case (3.5) reads

$$g_2^2 = \frac{5}{3} \left(\frac{1}{e^2} - \frac{1}{g_1^2} \right)^{-1} \quad (A.2)$$

which is then fixed by (3.4). The mass matrix M^2 in (2.20) is equal to

$$M^2 = g_0^T X_0 g = \frac{1}{2} (|a_1|^2 + |a_2|^2 + |a_3|^2) \begin{bmatrix} g_1^2 & -\sqrt{\frac{3}{5}} g_1 g_2 \\ -\sqrt{\frac{3}{5}} g_1 g_2 & \frac{3}{5} g_2^2 \end{bmatrix} = O'^T M_D^2 O' \quad (A.3)$$

$$O' = e \begin{bmatrix} g_1^{-1} & -\sqrt{\frac{5}{3}} g_2^{-1} \\ \sqrt{\frac{5}{3}} g_2^{-1} & g_1^{-1} \end{bmatrix}, \quad M_D^2 = \begin{bmatrix} 0 & 0 \\ 0 & M_1^2 \end{bmatrix},$$

with $|a_1|^2 + |a_2|^2 + |a_3|^2$ constrained by (3.6). Table 4 collects the corresponding fits.

(ii) The model with non-abelian flux-breaking requires to fix three columns namely

$$O_{ij} = \begin{bmatrix} \frac{\sqrt{3}}{2\sqrt{2}} & \frac{\sqrt{5}}{2\sqrt{2}} & 0 & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & -\frac{\sqrt{3}}{10} & \cdot \\ \frac{\sqrt{3}}{2\sqrt{2}} & -\frac{3}{2\sqrt{10}} & \frac{\sqrt{2}}{5} & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & -\frac{\sqrt{3}}{10} & \cdot \end{bmatrix} \quad (A.4)$$

It is worth noting that in this case the coupling of the extra current cannot be alternatively obtained by requiring large values of some of the vev's entering X in eq (2.13). One has to impose explicitly (project out) the decoupling of the very heavy gauge boson (which is heavy due to flux breaking). Eq (3.5) is then given by (A.2). For simplicity we assume in our fits in Table 5 that $g_3 = g_2$ given that the results are not appreciably affected by this choice (because of this we mark g_3 with an asterisk in Table 5). Finally

$$M^2 = g O^T X O g =$$

$$\frac{1}{2} \begin{bmatrix} g_1^2 (|a_1|^2 + |a_2|^2 + |a_3|^2), & -\frac{3}{5} g_1 g_2 (|a_1|^2 + |a_2|^2 + |a_3|^2), & -\frac{1}{\sqrt{15}} g_1 g_3 (|a_1|^2 + |a_2|^2 - 4|a_3|^2) \\ -\frac{3}{5} g_1 g_2 (|a_1|^2 + |a_2|^2 + |a_3|^2), & \frac{3}{5} g_2^2 (|a_1|^2 + |a_2|^2 + |a_3|^2), & \frac{1}{5} g_2 g_3 (|a_1|^2 + |a_2|^2 - 4|a_3|^2) \\ -\frac{1}{15} g_1 g_3 (|a_1|^2 + |a_2|^2 - 4|a_3|^2), & \frac{1}{5} g_2 g_3 (|a_1|^2 + |a_2|^2 - 4|a_3|^2), & \frac{1}{15} g_3^2 (|a_1|^2 + |a_2|^2 + 16|a_3|^2 + 25|a_4|^2 + 25|a_5|^2) \end{bmatrix}$$

Also $M^2 = O'^T M_D^2 O'$ where

$$O' = \begin{bmatrix} g_1^{-1} e & -\sqrt{\frac{5}{3}} g_2^{-1} e & 0 \\ \sqrt{\frac{5}{3}} g_2^{-1} e & g_1^{-1} e & 0 \\ 0 & 0 & 1 \end{bmatrix} \begin{bmatrix} 1 & 0 & 0 \\ 0 & \cos \delta_3 & \sin \delta_3 \\ 0 & -\sin \delta_3 & \cos \delta_3 \end{bmatrix}$$

(A.5)

$$M_D^2 = \begin{bmatrix} 0 & 0 & 0 \\ 0 & M_1^2 & 0 \\ 0 & 0 & M_2^2 \end{bmatrix},$$

Here we have explicitly decomposed $0'$ in a similar way to that in (3.1), (3.2). M_1 , M_2 and δ_3 depend on two parameters when (3.6) is used in (A.5), $(|a_1|^2 + |a_2|^2) / (|a_1|^2 + |a_2|^2 + |a_3|^2)$ and $|a_4|^2 + |a_5|^2$. Their explicit expressions follow immediately from (A.5).

(iii) The model of Table 6 is obtained by fixing

$$0_{ij} = \begin{bmatrix} \frac{\sqrt{3}}{2\sqrt{2}} & \frac{\sqrt{5}}{2\sqrt{2}} & 0 & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & -\frac{1}{2\sqrt{5}} & \cdot \\ \frac{\sqrt{3}}{2\sqrt{2}} & -\frac{3}{2\sqrt{10}} & -\frac{1}{2}\sqrt{\frac{3}{5}} & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{10}} & \frac{2}{\sqrt{5}} & \cdot \end{bmatrix} \quad (\text{A.6})$$

Eq. (A.2) applies and we assume $g_3 = g_2$ for this does not make a difference in the phenomenological fits (hence the asterisk on g_3 in Table 6). In this case the matrix is given by

$$M^2 = g_0^T X_0 g = \frac{1}{2} \begin{bmatrix} g_1^2 (|a_1|^2 + |a_2|^2 + |a_3|^2), -\sqrt{\frac{3}{5}} g_1 g_2 (|a_1|^2 + |a_2|^2 + |a_3|^2), -\frac{1}{\sqrt{10}} g_1 g_3 (2|a_1|^2 - 3|a_2|^2 + 2|a_3|^2) \\ -\sqrt{\frac{3}{5}} g_1 g_2 (|a_1|^2 + |a_2|^2 + |a_3|^2), \frac{3}{5} g_2^2 (|a_1|^2 + |a_2|^2 + |a_3|^2), \sqrt{\frac{3}{50}} g_2 g_3 (2|a_1|^2 - 3|a_2|^2 + 2|a_3|^2) \\ -\frac{1}{\sqrt{10}} g_1 g_3 (2|a_1|^2 - 3|a_2|^2 + 2|a_3|^2), \sqrt{\frac{3}{50}} g_2 g_3 (2|a_1|^2 - 3|a_2|^2 + 2|a_3|^2), \frac{1}{10} g_3^2 \times \\ \times (4|a_1|^2 + 9|a_2|^2 + 4|a_3|^2 + 25|a_5|^2) \end{bmatrix} \quad (\text{A.7})$$

$$= 0' M_D^2 0'^T,$$

with $0'$ and M_D^2 as in (A.5). M_1 , M_2 and δ_3 will depend on two parameters $(|a_1|^2 + |a_3|^2) / (|a_1|^2 + |a_2|^2 + |a_3|^2)$ and $|a_5|^2$ when (3.6) is used. Their dependence on these parameters follows from (A.7).

(iv) Table 7 corresponds to LR models.

$$O_{ij} = \begin{bmatrix} \frac{\sqrt{3}}{2\sqrt{2}} & \frac{\sqrt{3}}{2\sqrt{2}} & \frac{1}{2} & \cdot \\ \frac{1}{2\sqrt{2}} & \frac{1}{2\sqrt{2}} & -\frac{\sqrt{3}}{2} & \cdot \\ \frac{\sqrt{3}}{2\sqrt{2}} & -\frac{\sqrt{3}}{2\sqrt{2}} & 0 & \cdot \\ \frac{1}{2\sqrt{2}} & -\frac{1}{2\sqrt{2}} & 0 & \cdot \end{bmatrix} \quad (A.8)$$

$$g_2^2 = \left(\frac{1}{e^2} - \frac{1}{g_1^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-1}, \quad (A.9)$$

although in our fits we use $g_2 = g_1$ (because the fit is insensitive to this choice) and (A.9) fixes g_3 to be

$$g_3^2 = \frac{2}{3} \left(\frac{1}{e^2} - \frac{1}{g_1^2} - \frac{1}{g_2^2} \right)^{-1}. \quad (A.10)$$

Because of this we mark g_3 with an asterisk in Table 7.

M^2 is in this case equal to

$$g_0^T X_0 g =$$

$$\frac{1}{2} \begin{bmatrix} g_1^2 (|a_1|^2 + |a_2|^2 + |a_3|^2), -g_1 g_2 (|a_2|^2 + |a_3|^2), -\frac{3}{2} g_1 g_3 |a_1|^2 \\ -g_1 g_2 (|a_2|^2 + |a_3|^2), g_2^2 (|a_2|^2 + |a_3|^2 + |a_4|^2), -\frac{3}{2} g_2 g_3 |a_4|^2 \\ -\frac{3}{2} g_1 g_3 |a_1|^2, -\frac{3}{2} g_2 g_3 |a_4|^2, \frac{3}{2} g_3^2 (|a_1|^2 + |a_4|^2) \end{bmatrix} = O'^T M_D^2 O', \quad (A.11)$$

where O' is given by

$$O' = \begin{bmatrix} g_1^{-1} e , & -g_2^{-1} \left(\frac{1}{e^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-\frac{1}{2}} , & -\sqrt{\frac{2}{3}} g_1^{-1} g_3^{-1} e \left(\frac{1}{e^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-\frac{1}{2}} \\ g_2^{-1} e , & g_1^{-1} \left(\frac{1}{e^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-\frac{1}{2}} , & -\sqrt{\frac{2}{3}} g_2^{-1} g_3^{-1} e \left(\frac{1}{e^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-\frac{1}{2}} \\ \sqrt{\frac{2}{3}} g_3^{-1} e , & 0 , & e \left(\frac{1}{e^2} - \frac{2}{3} \frac{1}{g_3^2} \right)^{-\frac{1}{2}} \end{bmatrix} \begin{bmatrix} 1 & 0 & 0 \\ 0 & \cos \delta_3 & \sin \delta_3 \\ 0 & -\sin \delta_3 & \cos \delta_3 \end{bmatrix}$$

and M_D^2 as in (A.5). M_1, M_2 and δ_3 depend on $|a_1|^2 / (|a_1|^2 + |a_2|^2 + |a_3|^2)$ and $|a_4|^2$ when (3.6) is imposed. They can be obtained from (A.11).

(v) Finally Table 8 gives the fits for the general intermediate scale case. A large v.e.v. is assumed along a singlet direction, $|a_5|^2 \rightarrow \infty$. With no further assumptions the resulting model is similar to the LR one but with an extra parameter θ_3 .

$$O_{i\alpha} = O_{i\beta}^{LR} O_{\beta\alpha}^{(3)}, \quad i=1, \dots, 4, \beta, \alpha = 1, 2, 3, \text{ with } O^{LR} \text{ given by (A.8) and}$$

$$O_{\beta\alpha}^{(3)} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & \cos \theta_3 & \sin \theta_3 \\ 0 & -\sin \theta_3 & \cos \theta_3 \end{bmatrix} \quad (\text{A.12})$$

Then,

$$g_2^2 = (\cos \theta_3 - \sqrt{\frac{2}{3}} \sin \theta_3)^2 / \left[\frac{1}{e^2} - \frac{1}{g_1^2} - (\sin \theta_2 + \sqrt{\frac{2}{3}} \cos \theta_3)^2 \frac{1}{g_3^2} \right], \quad (\text{A.13})$$

and

$$M_{\alpha\beta}^2 = g_\alpha O_{\alpha\gamma}^{(3)T} M_{\gamma\delta}^2 O_{\delta\beta}^{(3)} g_\beta = O_{\alpha\gamma}^T M_{D\gamma\gamma}^2 O_{\gamma\beta}^T,$$

where M'^2 is given by

$$M'^2 = \frac{1}{2} \begin{pmatrix} |a_1|^2 + |a_2|^2 + |a_3|^2 , & -(|a_2|^2 + |a_3|^2) , & -\sqrt{\frac{3}{2}} |a_1|^2 \\ -(|a_2|^2 + |a_3|^2) , & |a_2|^2 + |a_3|^2 + |a_4|^2 , & -\sqrt{\frac{3}{2}} |a_4|^2 \\ -\sqrt{\frac{3}{2}} |a_1|^2 , & -\sqrt{\frac{3}{2}} |a_4|^2 , & \frac{3}{2} (|a_1|^2 + |a_4|^2) \end{pmatrix}, \quad (\text{A.14})$$

M_D^2 is given by (A.5) and

$$O' = \begin{pmatrix} c_1 c_2 & s_1 & -s_2 c_1 \\ -s_1 c_2 & c_1 & s_1 s_2 \\ s_2 & 0 & c_2 \end{pmatrix} \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \delta_3 & \sin \delta_3 \\ 0 & -\sin \delta_3 & \cos \delta_3 \end{pmatrix},$$

with $s_2 = g_3^{-1} e (\sin \theta_3 + \sqrt{\frac{2}{3}} \cos \theta_3)$, $c_2 = e \left(\frac{1}{e^2} - \frac{1}{g_3^2} (\sin \theta_3 + \sqrt{\frac{2}{3}} \cos \theta_3)^2 \right)^{\frac{1}{2}}$,
 $s_1 = -g_2^{-1} e (\cos \theta_2 - \sqrt{\frac{2}{3}} \sin \theta_2) / c_2$, $c_1 = g_1^{-1} e / c_2$. As in the former cases
 after imposing (3.6) M_1 , M_2 and δ_3 depend on $|a_1|^2 / (|a_1|^2 + |a_2|^2 + |a_3|^2)$ and
 $|a_4|^2$ (besides θ_3 and g_1).

Table Captions

Table 1

Generators of $U_1^{T^L} x U_1^{Y^L} x U_1^{T^R} x U_1^{Y^R}$ acting on the 27 of E_6 .

Table 2

Phenomenological coefficients describing neutral current interactions. The given expressions are derived on the assumption that the observed fermions correspond to the states with subscript 1 in Table 1.

Table 3

Neutral current data fits to (E_6) superstring inspired models with general diagonal gauge interactions. The first line corresponds to $\chi^2/\text{number of degrees of freedom (data points-free parameters)}$. $\sin^2 \theta_W = \frac{e^2}{g_1^2}$ is fixed in the first two columns to reproduce the measured UA1 and UA2 values for M_W . In the third column the values of M_W are used as data points. $\theta_{1,3} \in [-\pi, \pi]$, whereas $\theta_2 \in [-\frac{\pi}{2}, \frac{\pi}{2}]$, similarly for $\delta_{1,3}$ and δ_2 . (Note that not all the parameter space can be distinguished for the signs of currents and gauge bosons are not relevant). x, y, u and v are positive quantities defined as follows: $|a_1|^2 = x(2/2 G_F)^{-1}$, $|a_2|^2 = y(1-x)(2/2 G_F)^{-1}$, $|a_3|^2 = (1-y)(1-x)(2/2 G_F)^{-1}$, $|a_4|^2 = u(2/2 G_F)^{-1}$, $|a_5|^2 = v(2/2 G_F)^{-1}$. These quantities are quoted with parabolic errors. $M_Z = M_1$ in the text. The first subcolumns correspond to the case of symmetry breaking due to $27(\overline{27})$'s, and the second ones correspond to fits with arbitrary gauge boson masses and mixings. Only free parameters are quoted with errors. We discuss them in the text.

Table 4

Fits of neutral current data to the standard model. See Table 3 and Appendix for details.

Table 5

Fits of neutral current data to the non-abelian flux breaking model. See Table 3 and Appendix for details. Here x and u are given by $|a_1|^2 + |a_2|^2 = x(2/2 G_F)^{-1}$, $|a_3|^2 = (1-x)(2/2 G_F)^{-1}$, $|a_4|^2 + |a_5|^2 = u(2/2 G_F)^{-1}$. Errors are parabolic, but we also include for the mass of the second gauge boson the correct non-parabolic lower limit (underlined).

Table 6

Fits of neutral current data to an intermediate scale model with a v.e.v. along the v_R direction. See Table 3 and Appendix for details. Here x and u are given by $|a_1|^2 + |a_3|^2 = x(2/2 G_F)^{-1}$, $|a_2|^2 = (1-x)(2/2 G_F)^{-1}$, $|a_5|^2 = u(2/2 G_F)^{-1}$. Errors are parabolic, but we also include for the mass of the second gauge boson the correct non-parabolic lower limit (underlined).

Table 7

Fits of neutral current data to a left-right symmetric model. See Table 3 and Appendix for details. Here x and u are defined by $|a_1|^2 = x(2/2 G_F)^{-1}$, $|a_2|^2 + |a_3|^2 = (1-x)(2/2 G_F)^{-1}$, $|a_4|^2 = u(2/2 G_F)^{-1}$. Errors are parabolic, but we include also for the mass of the second gauge boson the correct non-parabolic lower limit (underlined).

Table 8

Fits of neutral current data to the general intermediate scale model. See Table 3 and Appendix for details. The quantities x and u are defined as in Table 7.

Table 1

	T_3^L	Y^L	T_3^R	Y^R
$\begin{pmatrix} u \\ d_1 \end{pmatrix}_L$	$\frac{1}{2}$ $-\frac{1}{2}$	$\frac{1}{2\sqrt{3}}$ $\frac{1}{2\sqrt{3}}$	0 0	0 0
d_{2L}	0	$-\frac{1}{\sqrt{3}}$	0	0
u_L^c	0	0	$-\frac{1}{2}$	$-\frac{1}{2\sqrt{3}}$
d_{1L}^c	0	0	$\frac{1}{2}$	$-\frac{1}{2\sqrt{3}}$
d_{2L}^c	0	0	0	$\frac{1}{\sqrt{3}}$
$\begin{pmatrix} \nu_1 \\ e_1 \end{pmatrix}_L$	$\frac{1}{2}$ $-\frac{1}{2}$	$-\frac{1}{2\sqrt{3}}$ $-\frac{1}{2\sqrt{3}}$	0 0	$-\frac{1}{\sqrt{3}}$ $-\frac{1}{\sqrt{3}}$
$\begin{pmatrix} \nu_2 \\ e_2 \end{pmatrix}_L$	$\frac{1}{2}$ $-\frac{1}{2}$	$-\frac{1}{2\sqrt{3}}$ $-\frac{1}{2\sqrt{3}}$	$-\frac{1}{2}$ $-\frac{1}{2}$	$\frac{1}{2\sqrt{3}}$ $\frac{1}{2\sqrt{3}}$
$\begin{pmatrix} e_2^c \\ \nu_3 \end{pmatrix}_L$	$\frac{1}{2}$ $-\frac{1}{2}$	$-\frac{1}{2\sqrt{3}}$ $-\frac{1}{2\sqrt{3}}$	$\frac{1}{2}$ $\frac{1}{2}$	$\frac{1}{2\sqrt{3}}$ $\frac{1}{2\sqrt{3}}$
e_{1L}^c	0	$\frac{1}{\sqrt{3}}$	$\frac{1}{2}$	$\frac{1}{2\sqrt{3}}$
ν_{4L}	0	$\frac{1}{\sqrt{3}}$	$-\frac{1}{2}$	$\frac{1}{2\sqrt{3}}$
ν_{5L}	0	$\frac{1}{\sqrt{3}}$	0	$-\frac{1}{\sqrt{3}}$

Table 2

$$\frac{4G_F}{\sqrt{2}} \alpha = -\frac{1}{2} (\sqrt{3} B_{\tilde{Q}A} + B_{\tilde{Q}B} + B_{AA} + \frac{1}{\sqrt{3}} B_{AB})$$

$$\frac{4G_F}{\sqrt{2}} \beta = -\frac{1}{2} (\sqrt{3} B_{AB} + B_{AC} + B_{BB} + \frac{1}{\sqrt{3}} B_{BC})$$

$$\frac{4G_F}{\sqrt{2}} \gamma = -\frac{1}{2} (\frac{1}{\sqrt{3}} B_{\tilde{Q}A} + \frac{1}{3} B_{\tilde{Q}B} - B_{AA} - \frac{1}{\sqrt{3}} B_{AB})$$

$$\frac{4G_F}{\sqrt{2}} \delta = -\frac{1}{2} (\frac{1}{\sqrt{3}} B_{AB} - B_{AC} + \frac{1}{3} B_{BB} - \frac{1}{\sqrt{3}} B_{BC})$$

$$\frac{4G_F}{\sqrt{2}} g_V = \frac{1}{2} (\sqrt{3} B_{\tilde{Q}A} + B_{\tilde{Q}B} - B_{AA} - \frac{1}{\sqrt{3}} B_{AB})$$

$$\frac{4G_F}{\sqrt{2}} g_A = \frac{1}{2} (\frac{1}{\sqrt{3}} B_{AB} + B_{AC} + \frac{1}{3} B_{BB} + \frac{1}{\sqrt{3}} B_{BC})$$

$$\frac{4G_F}{\sqrt{2}} \tilde{\alpha} = \frac{1}{2} (B_{\tilde{Q}B} + \sqrt{3} B_{\tilde{Q}C} + \frac{1}{\sqrt{3}} B_{AB} + B_{AC})$$

$$\frac{4G_F}{\sqrt{2}} \tilde{\beta} = \frac{1}{2} (3B_{\tilde{Q}B} + \sqrt{3} B_{\tilde{Q}C} - \sqrt{3} B_{AB} - B_{AC})$$

$$\frac{4G_F}{\sqrt{2}} \tilde{\gamma} = \frac{1}{2} (\frac{1}{3} B_{\tilde{Q}B} + \frac{1}{\sqrt{3}} B_{\tilde{Q}C} - \frac{1}{\sqrt{3}} B_{AB} - B_{AC})$$

$$\frac{4G_F}{\sqrt{2}} \tilde{\delta} = \frac{1}{2} (B_{\tilde{Q}B} - \sqrt{3} B_{\tilde{Q}C} - \frac{1}{\sqrt{3}} B_{AB} + B_{AC})$$

$$\frac{4G_F}{\sqrt{2}} h_{AA}^{\mu,b} = -\frac{1}{4} (\frac{1}{3} B_{BB} + \frac{2}{\sqrt{3}} B_{BC} + B_{CC})$$

$$\frac{4G_F}{\sqrt{2}} h_{AA}^c = \frac{1}{2} (\frac{1}{3} B_{BB} + \frac{1}{\sqrt{3}} B_{BC})$$

$$B_{ij} = B_{ji} = (0g0' [q^2 - M_D^2]^{-1} 0'^T g0^T)_{ij} = [0g^{-2} 0^T q^2 - x]_{ij}^{-1}, \quad i, j = \tilde{Q}, A, B, C$$

Table 3

	$M_W = 83.5$		$M_W = 81.2$			
$\chi^2 / \#$	66.8/75 35.1/42	62.9/73 29.3/40	64.7/75 29.3/42	63.7/73 29.1/40	66.1/76 29.6/43	63.5/74 29.4/41
$\sin^2 \theta_W$	.21 .21	.21 .21	.23 .23	.23 .23	.22±.01 .22±.01	.22±.01 .22±.01
g_3	.42±.11 .34±.00	.38±.15 .79±.78	.44±.10 .19±.59	.25±.24 .55±.19	.50±.34 .16±.06	.32±.13 .50±.28
g_4	1.00±.64 .46±.18	.83±.13 .77±.14	1.00±.71 .80±.83	1.00±.19 .74±.32	.79±.00 .79±.83	.90±.23 .92±.13
θ_1	3.07±.40 2.28±.49	-2.69±.13 -2.94±.32	-2.71±4.66 2.89±1.21	-2.82±.12 -2.73±.17	2.86±.40 3.11±.39	-2.71±.12 -2.82±.11
θ_2	-1.57±.14 -1.25±.46	-1.28±.32 -1.45±.12	-1.34±.60 -1.48±3.06	-1.06±.37 -1.17±.18	-1.55±.13 -1.56±2.21	-1.23±.24 -1.18±.93
θ_3	2.13±.77 2.34±.53	2.32±.08 2.22±.06	2.20±4.92 2.24±.34	2.37±.19 2.15±.17	-1.35±.36 2.36±.19	2.33±.11 2.19±.20
x	.30±.18 .64±.31		.34±.13 .00±.71		.25±.08 .00±1.00	
y	1.00±.00 .25±.59		1.00±.84 .72±.79		1.00±.75 .86±.29	
u	15.2±4.8 26.9±62.1		12.1±13.3 75.6±82.5		7.8±4.1 69.4±79.3	
v	27.1±.0 16.4±9.6		8.4±7.1 15.9±30.4		52.0±74.2 25.8±63.4	
M_Z	92.4 94.0	92.6±2.1 92.4±1.9	92.0 92.1	92.0±1.7 92.3±1.8	92.2 92.3	92.4±2.2 92.3±1.1
M_2	265 439	150±44 163±55	207 246	160±82 148±134	226 255	179±70 227±75
M_3	425 2329.5	204±71 942±956	521 442	286±105 451±359	469 313	184±45 540±768
δ_1	.69 -1.57	-1.55±.80 -1.57±.72	.29 .55	.00±.95 -1.52±.60	-.43 .41	-1.57±.41 -1.57±4.04
δ_2	.16 .24	.01±.17 .11±.09	.15 .26	.12±.05 .11±.03	-.27 .27	.06±.07 .11±.06
δ_3	3.12 -1.73	2.83±.19 3.04±.14	2.89 -2.85	2.85±.10 2.92±.04	-3.04 -2.97	2.78±.15 2.95±.10

Table 4

	$M_W = 83.5$		$M_W = 81.2$			
$\chi^2 / \#$	75.8/84 37.2/51	72.2/83 35.2/50	73.4/84 31.6/51	72.4/83 31.5/50	73.6/85 31.8/52	71.9/84 31.2/51
$\sin^2 \theta_W$	.21 .21	.21 .21	.23 .23	.23 .23	.22±.01 .23±.01	.22±.01 .23±.01
M_Z	94.1 94.1	95.0±.4 94.8±.4	92.3 92.3	92.7±.4 92.4±.4	92.6 91.9	93.7±1.2 91.6±1.6

Table 5

	$M_W = 83.5$		$M_W = 81.2$			
$\chi^2 / \#$	75.9/82 37.2/49	72.5/81 35.3/48	73.4/82 31.6/49	72.7/81 31.5/48	73.6/83 31.8/50	72.3/82 31.8/49
$\sin^2 \theta_W$	.21 .21	.21 .21	.23 .23	.23 .23	.22±.01 .23±.01	.22±.01 .23±.01
g_3^*	.46 .46	.46 .46	.46 .46	.46 .46	.46 .46	.46 .46
x	1.00±.92 .92±.94		1.00±.74 .96±.52		1.00±.82 1.00±1.00	
u	101.0±86.4 101.0±94.0		101.0±96.5 101.0±50.7		101.0±50.7 101.0±53.2	
M_Z	94.1 94.1	95.0±.4 94.8±.5	92.3 92.3	92.7±.4 92.4±.4	92.6 91.9	93.7±1.2 92.1±1.3
M_2	728 750 ± 645 $\underline{-383}$		734 750 ± 638 $\underline{-370}$		733 750 ± 642 $\underline{-376}$	
	729 750 ± 402 $\underline{-527}$		734 750 ± 601 $\underline{-555}$		735 750 ± 598 $\underline{-581}$	
δ_3	.003 .002	.000±.009 .000±.060	.003 .002	.000±.007 .000±.113	.003 .003	.000±.008 .000±.118

Table 6

	$M_W = 83.5$		$M_W = 81.2$			
$\chi^2 / \#$	72.3/82	67.2/81	69.5/82	66.3/81	69.8/83	66.3/82
	34.8/49	33.7/48	30.9/49	30.8/48	31.2/50	31.1/49
$\sin^2 \theta_W$	.21	.21	.23	.23	.22 $\pm$.01	.22 $\pm$.01
	.21	.21	.23	.23	.23 $\pm$.01	.23 $\pm$.01
g_3^*	.46	.46	.46	.46	.46	.46
	.46	.46	.46	.46	.46	.46
x	.61 $\pm$.08		.56 $\pm$.07		.57 $\pm$.08	
	.66 $\pm$.11		.61 $\pm$.17		.60 $\pm$.17	
u	4.4 $\pm$ 2.5		4.0 $\pm$ 2.2		4.1 $\pm$ 2.3	
	3.9 $\pm$ 2.8		6.3 $\pm$ 7.9		6.6 $\pm$ 8.7	
M_Z	94.1	95.5 $\pm$.5	92.3	93.2 $\pm$.5	92.5	93.9 $\pm$ 1.3
	94.1	94.9 $\pm$.7	92.3	92.5 $\pm$.7	92.1	92.4 $\pm$ 1.5
M_2	192	192 ± 56 <u>-40</u>	185	186 ± 49 <u>-36</u>	186	188 ± 53 <u>-37</u>
	181	181 ± 73 <u>-47</u>	228	228 ± 135 <u>-76</u>	234	230 ± 147 <u>-80</u>
δ_3	.004	-.090 $\pm$.061	-.013	-.091 $\pm$.057	-.011	-.090 $\pm$.060
	.023	-.022 $\pm$.063	.002	-.008 $\pm$.084	.001	-.008 $\pm$.108

Table 7

	$M_W = 83.5$		$M_W = 81.2$							
$\chi^2/\#$	72.9/82 36.1/49	70.1/81 32.9/48	71.4/82 31.6/49	71.1/81 31.1/48	71.4/83 31.9/50	70.6/82 31.4/49				
$\sin^2\theta_W$	.21 .21	.21 .21	.23 .23	.23 .23	.22±.01 .23±.01	.22±.01 .22±.01				
g_3^*	.34 .34	.34 .34	.35 .35	.35 .35	.35 .35	.34 .34				
x	.80±.19 .90±.21		.73±.30 1.00±.96		.76±.26 1.00±.51					
u	13.7±8.1 16.5±16.7		15.8±12.0 101.0±90.6		15.4±9.8 101.0±53.9					
M_Z	94.1 94.1	94.7±.5 94.4±.5	92.3 92.2	92.5±.5 92.3±.5	92.8 91.9	93.6±1.3 92.4±1.4				
M_2	364 398	491 1000	±298 -172	385 970	411 1000	±187 -116	382 968	448 1000	±243 -143	±973 -460
δ_3	-.644 -.637	-.627±.036 -.595±.033	-.636 -.635	-.639±.037 -.613±.034	-.637 -.634	-.633±.037 -.611±.035				

Table 8

	$M_W = 83.5$		$M_W = 81.2$			
$\chi^2 / \#$	68.0/81 30.9/48	67.8/80 29.6/47	66.5/81 29.7/48	66.5/80 29.6/47	66.8/82 30.0/49	66.7/81 29.8/48
$\sin^2 \theta_W$	.21 .21	.21 .21	.23 .23	.23 .23	.22±.01 .22±.01	.22±.01 .22±.00
g_3	.50±.01 .97±.70	.63±.29 .59±.02	.77±.20 .80±.31	.59±.07 .58±.12	.53±.62 .83±.23	.59±.07 .55±.01
θ_3	.72±.03 .54±.08	.37±.56 .78±.09	-.02±.15 .43±.22	.20±.18 .67±.26	.48±4.86 .44±.21	.22±.21 .75±.10
x	.16±.07 .00±.12		.24±.05 .00±.78		.24±.26 .00±.69	
u	8.8±2.4 21.0±13.9		6.2±1.9 19.5±19.2		6.4±6.9 19.5±20.2	
M_Z	92.6 93.2	94.5±1.7 92.6±2.3	92.2 92.1	92.5±1.1 92.0±1.4	92.4 92.3	92.9±1.7 92.4±.5
M_2	110 333	218±182 141±18	320 357	223±46 194±132	173 352	223±53 147±15
δ_3	-.461 -.182	-.331±.161 -.329±.367	-.175 -.259	-.336±.018 -.492±.263	-.536 -.248	-.340±.020 -.531±.022

References

- [1] J. Schwarz, Phys. Rep. 89 (1982) 223; M. Green and J. Schwarz, Phys. Lett. 149B (1984) 117; D.T. Gross, J.A. Harvey, E. Martinec and R. Rohm, Phys. Lett. 54 (1985) 502; Nucl. Phys. B256 (1985) 253; Nucl. Phys. B267 (1986) 75.
- [2] P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, Nucl. Phys. B258 (1985) 46; E. Witten, Nucl. Phys. B258 (1985) 75; M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. B259 (1985) 549; J.D. Breit, B. Ovrut and G. Segre, Phys. Lett. 158B (1985) 33; S. Cecotti, J.P. Derendinger, S. Ferrara, L. Girardello and M. Roncadelli, Phys. Lett. 156B (1985) 318; E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 161B (1985) 85; Phys. Lett. 165B (1985) 76; J.P. Derendinger, L.E. Ibanez and H.P. Nilles, Phys. Lett. 155B (1985) 65; Nucl. Phys. B267 (1986) 365.
- [3] I. Gursev, P. Ramond and P. Sikivie, Phys. Lett. 60B (1976) 177; Y. Achiman and B. Stech, Phys. Lett. 77B (1978) 389; Q. Shafi, Phys. Lett. 79B (1978) 301; H. Ruegg and T. Schuler, Nucl. Phys. B161 (1979) 388. R. Barbieri and D.V. Nanopoulos Phys. Lett. 91B (1980) 369.
- [4] J.E. Kim, P. Langacker, M. Levine and M.M. Williams, Rev. Mod. Phys. 53 (1981) 211, and references therein.
- [5] G. Barbiellini and C. Santoni, CERN preprint - EP/85-117, submitted to Riv. Nuovo Cimento and references therein.
- [6] C.Y. Prescott et al., Phys. Lett. 84B (1979) 524.
- [7] A. Argento et al., Phys. Lett. 120B (1983) 245.
- [8] L. DiLella, Physics from the $p\bar{p}$ collider, CERN preprint-EP/85-184, Rapporteur's talk at Bari, 18-24 July 1985, and references therein.
- [9] P.Q. Hung and J.J. Sakurai, Phys. Lett. 69B (1977) 323.
- [10] H. Georgi and S. Weinberg, Phys. Rev. D17 (1978) 275; F. del Aguila and A. Mendez, Nucl. Phys. B189 (1981) 212.
- [11] J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, CERN preprint -TH4323/85; V. Barger, N.G. Deshpande and K. Whisnant, Phys. Rev. Lett. 56 (1986) 30; see also L.S. Durkin and P. Langacker, Phys. Lett. 166B (1986) 436.
- [12] L. Ibanez, CERN preprint TH-4308/85, to be published in the Proceedings of the 1st Torino Meeting, World Scientific Pub. Co.; L. Ibanez and J. Mas, CERN preprint TH-4426.

- [13] A. De Rujula, H. Georgi and S.L. Glashow, Ann. Physics 109 (1977) 242; 258; R.N. Mohapatra and D.P. Sidhu, Phys. Rev. Lett. 38 (1977) 665.
- [14] S.L. Glashow, Nucl. Phys. 22 (1961) 579;
S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264;
A. Salam, Proc. 8th Nobel Symp., ed. N. Svartholm (Almqvist and Wiksells, Stockholm, 1968) p. 367.
- [15] F. del Aguila and M.J. Bowick, Nucl. Phys. B224 (1983) 107;
see also: P. Ramond, Proc. 4th Kyoto Summer Institute (Kyoto, Japan, 1981), eds. M. Kommuna and T. Maskawa (World Scientific, Singapore, 1981); P. Fishbane, S. Meshkov, R.E. Norton and P. Ramond, Phys. Rev. D31 (1985) 1119.
- [16] K. Choi and J.E. Kim, Preprint SNUHE 86/04.
- [17] E. Bellotti et al., Phys. Lett. 121B (1983) 72.
- [18] F.E. James and M. Roos, Comput. Phys. Commun. 10 (1975) 343.
- [19] J. Bernabéu and C. Jarlskog, Phys. Lett. 69B (1977) 71;
S.M. Barr and A. Zee, Phys. Lett. 92B (1980) 297, 110B(1982) 501;
S.M. Barr, Phys. Lett. 128B (1983) 400.
- [20] S.M. Barr, Phys. Rev. Lett. 55 (1985) 2778.
- [21] F. del Aguila, G. Blair, M. Daniel and G.G. Ross, Nucl. Phys. B272 (1986) 413.
- [22] N.G. Deshpande and R.L. Johnson, Phys. Rev. D27 (1983) 1165;
see also Ref [13].
- [23] C. Mana, UAM-Ph. Thesis (1985) and references therein.
- [24] V. Barger, N.G. Deshpande and K. Whisnant, Phys. Rev. Lett. 56 (1986) 30.
- [25] M.A. Bouchiat et al., Phys. Lett. 134B (1984) 463.
- [26] G. Fogli, Nucl. Phys. B207 (1982) 322.
- [27] R.W. Robinett, Phys. Rev. D33 (1986) 1908; V. Barger, N.G. Deshpande, R.J.N. Phillips and K. Whisnant, Phys. Rev. D33 (1986) 1912.
- [28] M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. B259 (1985) 549.