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DETECTING E_6 NEUTRAL GAUGE BOSONS THROUGH LEPTON PAIRS AT HADRON COLLIDERS

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ABSTRACT

We examine the possibility of observing new neutral gauge bosons from E_6 superstring-inspired models at hadron colliders. Cross-sections and asymmetries for Drell-Yan production of lepton-antilepton pairs are given. A general parametrization for the cases of one and two extra neutral gauge bosons is discussed and currently popular models are included as particular examples. Results are presented for the $p\bar{p}$ colliders at CERN and Fermilab and for the planned pp supercolliders, LHC and SSC.

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1. - INTRODUCTION

Recent theoretical work seems to indicate that Nature may be described by consistent theories of extended objects in higher dimensions, superstrings¹⁾. In particular, the heterotic superstring²⁾, formulated in 10 dimensions with gauge group $E_8 \times E_8$, seems to be particularly promising for phenomenology, thanks to a class of compactifications³⁾ to four dimensions giving rise to chiral fermions, unbroken $N = 1$ supersymmetry and a subgroup of E_6 ⁴⁾ as the effective low-energy group⁵⁾. These compactifications make plausible the existence of new gauge bosons at mass scales that can be tested by present or planned accelerators. In this paper, we examine the possibility of observing new gauge bosons, associated to E_6 (diagonal) neutral currents, at hadron colliders. Although there is an extensive literature on the existence of new gauge bosons⁶⁾⁻¹⁰⁾, only particular cases are usually discussed. We analyze the general E_6 case with one and two extra gauge bosons in the TeV range. Popular models^{9),11),12)} with one extra gauge boson are treated as particular cases.

Considering only decays into known particles, the most efficient way of observing new gauge bosons in proton-(anti)proton colliders is to identify charged leptons in the final state. Here we concentrate^{*)} on Drell-Yan production of lepton-antilepton pairs, $p(\bar{p}) \rightarrow \ell^+ \ell^- X$ [$\ell = e, \mu, \dots$], which allowed for the discovery of the Z at the CERN $p\bar{p}$ -collider¹⁴⁾. Other signals of extra neutral gauge bosons, as well as possible backgrounds, will be discussed elsewhere. Decays into exotic fermions or into supersymmetric particles are only considered here for their contribution to the width of the new gauge boson. Constraints on its mass and its cross-section into W-pairs were given in Ref. 15).

The unpolarized differential cross-section for the process $p(\bar{p}) \rightarrow \ell^+ \ell^- X$ depends on the lepton invariant mass M , on the rapidity y and on the angle θ^* in the centre of mass of the colliding partons (see Fig. 1). The general form of this cross-section is, for tree level amplitudes,

*) We thank D.V. Nanopoulos for urging us to study the possibility of observing new gauge bosons at the SSC pp supercollider, which motivated this work. A dedicated study for the SSC can be found in Ref. 13).

$$\frac{d\sigma}{dM dy d\cos\theta^*} = \sum_{q=\text{quarks}} \left[g_q^S(y, M) S_q(M) (1 + \cos^2\theta^*) + g_q^A(y, M) A_q(M) 2 \cos\theta^* \right], \quad (1.1)$$

where S_q , A_q (which are the only model-dependent quantities) involve the vector and axial couplings of quarks and leptons to the different neutral gauge bosons, while g_q^S , g_q^A involve the parton distribution functions of the colliding hadrons: the precise definitions will be given in Section 2. One usually integrates the θ^* dependence, which only reflects the spin structure of the interaction. One can also integrate the y -dependence, which gives the parton distribution within the hadrons and is an input, or the M dependence, given that almost all the events concentrate on the resonant peak. However, one loses in principle a lot of information in both cases. If the rapidity dependence is integrated out, given that the structure functions weigh differently the (anti)quark distributions for different y values, one gets only one average of S_q and A_q , while keeping explicitly the y -dependence one might probe different combinations of S_q and A_q . On the other hand, if the M -dependence is integrated out, one loses in a similar way relevant interference effects, as S_q and A_q are sums of products of coupling constants and propagators. If a signal is detected, the best way of analyzing the data would be to perform a fit to (1.1), given that the θ^* and y -dependences for the production of a new Z are calculable. For this the general parametrization of the E_6 models we consider is convenient, because the preferred model can be singled out at once. However, the only sensible thing to do at present (in the absence of any experimental indication of a new Z) is to give numerical examples of differential cross-sections and asymmetries, integrating out θ^* and y or θ^* and M , respectively.

Experimentally, the situation can be more complicated. To discriminate among different models, the lepton charges have to be identified and the momenta measured with a sufficient resolution: this could be non-trivial at LHC and SSC¹⁶⁾. Moreover, to take cuts and detection efficiency into account, one should multiply the lepton cross-section by an experiment-dependent function of M , y and θ^* .

Finally, the results depend on the parton distribution functions: in this paper we use those of Duke and Owens¹⁷⁾, set 1; other choices^{17), 18)} would not affect significantly the results of our analysis. QCD corrections can also be

non-negligible¹⁹⁾, and in particular partons can carry a non-vanishing transverse momentum. Here we neglect all these effects.

In Section 2, we calculate the differential cross-section (1.1) for the general E_6 case. Details of the parametrization of the different models are given in the Appendices. We discuss the possibility of one or two extra neutral gauge bosons. Popular models^{9),11),12)} with one extra gauge boson are treated as particular cases.

In Section 3, we plot various differential cross-sections and asymmetries, as well as the total cross-section, for different ranges of gauge boson masses and coupling constants. We concentrate on the case of one extra gauge boson. Figures are presented for the $p\bar{p}$ colliders at CERN and Fermilab and for the planned pp supercolliders (LHC and SSC).

In Section 4 we draw our conclusions and make some additional comments.

2. - LEPTON DIFFERENTIAL CROSS-SECTIONS AND ASYMMETRIES

The production and decay of neutral electroweak gauge bosons is described at the lowest order by the general Lagrangian

$$\mathcal{L}_{NC} = J_{\alpha\mu} Z_{\alpha}^{\mu} = \bar{\psi}^k \gamma_{\mu} (v_{\alpha}^k + a_{\alpha}^k \gamma_5) \psi^k Z_{\alpha}^{\mu}, \quad (2.1)$$

where summation over repeated indices is understood. Z_{α} ($\alpha = 1, 2, 3, \dots$) stands for the physical gauge bosons of definite mass M_{α} and width Γ_{α} : $\alpha = 1$ corresponds to the photon ($M_1 = 0$, $\Gamma_1 = 0$), $\alpha = 2$ to the observed Z ($M_2 = 92.6 \pm 1.7$ GeV, $\Gamma_2 < 4.6$ GeV)^{14),20)}, $\alpha = 3, 4$ to possible extra neutral gauge bosons. The index k runs over the different fermions. For the standard model $\alpha = 1, 2$, and the corresponding coefficients v_{α}^k , a_{α}^k for the observed quarks and leptons are given in Table 1.

At hadron colliders, the production of lepton-antilepton pairs is described, at the quark level, by the graphs of Fig. 2. Searching for isolated lepton-antilepton pairs with large transverse energy one can observe the Z and possibly new neutral gauge bosons. The corresponding differential cross-section is given by Eq. (1.1), with

$$g_q^{S,A}(y, M) = \frac{M}{48\pi} x_a x_b \left[f_q^{(a)}(x_a, M^2) f_{\bar{q}}^{(b)}(x_b, M^2) \pm f_{\bar{q}}^{(a)}(x_a, M^2) f_q^{(b)}(x_b, M^2) \right], \quad (2.2)$$

where the sign $+(-)$ corresponds to $S(A)$, and

$$S_q = \sum_{\alpha, \beta} \frac{(v_\alpha^q v_\beta^q + a_\alpha^q a_\beta^q)(v_\alpha^l v_\beta^l + a_\alpha^l a_\beta^l)}{(M^2 - M_\alpha^2 + i M_\alpha \Gamma_\alpha)(M^2 - M_\beta^2 - i M_\beta \Gamma_\beta)}, \quad (2.3)$$

$$A_q = \sum_{\alpha, \beta} \frac{(v_\alpha^q a_\beta^q + v_\beta^q a_\alpha^q)(v_\alpha^l a_\beta^l + v_\beta^l a_\alpha^l)}{(M^2 - M_\alpha^2 + i M_\alpha \Gamma_\alpha)(M^2 - M_\beta^2 - i M_\beta \Gamma_\beta)}.$$

We follow standard conventions²¹⁾ and the notation of Ref. 18) if not otherwise stated. Quark and lepton masses are neglected. a and b are the two colliding hadrons at centre-of-mass energy $\sqrt{s}$, x_a and x_b being the momentum fractions of the colliding partons in a and b respectively (see Fig. 1). M is the lepton-antilepton invariant mass, y the rapidity ($x_{a,b} = M/\sqrt{s} e^{\pm y}$) and θ^* the scattering angle ($a\bar{l}^-$) in the centre of mass of the parton system $q\bar{q}$. Finally, $f_{q(\bar{q})}^{(a(b))}(x_{a(b)}, M^2)$ are the relevant parton distribution functions.

All the model dependence in Eq. (1.1) is in S_q and A_q . Then it is sensible to integrate over θ^* . We will present plots of the differential cross-section

$$\left(\frac{d\sigma}{dM} \right)^S = \int_{\ln \frac{M}{\sqrt{s}}}^{\ln \frac{\sqrt{s}}{M}} dy \int_{-1}^1 d\cos\theta^* \frac{d\sigma}{dM dy d\cos\theta^*}, \quad (2.4)$$

testing the dependence of S_q on the model parameters, and the asymmetry for $p\bar{p}$

$$A^{p\bar{p}}(M) = \frac{(d\sigma/dM)^{A^{p\bar{p}}}}{(d\sigma/dM)^S} =$$

$$= \frac{\int_{\ln M/\sqrt{s}}^{\ln \sqrt{s}/M} dy \left(\int_0^1 - \int_{-1}^0 \right) d\cos\theta^* \frac{d\sigma}{dM dy d\cos\theta^*}}{(\frac{d\sigma}{dM})^S}, \quad (2.5a)$$

and for pp

$$A^{pp}(M) = \frac{(\frac{d\sigma}{dM})^{APP}}{(\frac{d\sigma}{dM})^S} = \frac{\left(\int_0^{\ln \sqrt{s}/M} - \int_{\ln M/\sqrt{s}}^0 \right) dy \left(\int_0^1 - \int_{-1}^0 \right) d\cos\theta^* \frac{d\sigma}{dM dy d\cos\theta^*}}{(\frac{d\sigma}{dM})^S}, \quad (2.5b)$$

testing the A_q dependence (for a sizeable differential cross-section). We also give the differential cross-section $(\frac{d\sigma}{dM dy})_{M=M_3}^S$ at the resonant peak $M = M_3$ and the corresponding asymmetry

$$A(y) = \frac{(\frac{d\sigma}{dM dy})_{M=M_3}^A}{(\frac{d\sigma}{dM dy})_{M=M_3}^S} = \frac{\left(\int_0^1 - \int_{-1}^0 \right) d\cos\theta^* \frac{d\sigma}{dM dy d\cos\theta^*} \Big|_{M=M_3}}{\int_{-1}^1 d\cos\theta^* \frac{d\sigma}{dM dy d\cos\theta^*} \Big|_{M=M_3}}. \quad (2.6)$$

Equations (2.4)-(2.6) are our main concern in the next section. However, let us first discuss how S_q and A_q depend on the E_6 model parameters.

In order to fix the relevant Lagrangian (2.1), we have to discuss the gauge boson mass generation and the fermion assignments within the E_6 fundamental representation, 27. We first discuss the parametrization required by the gauge boson mass generation. E_6 has four diagonal $SU(3)_C \times U(1)_Q$ neutral gauge bosons Z_α , that can be identified with the photon, the observed Z and two new ones. If for some mechanism the last two acquire large masses, the low-energy Lagrangian in (2.1) reduces to the standard model: $\alpha = 1, 2$, $M_Y = \Gamma_Y = 0$, $M_{Z_2} \sim 92.6$ GeV,

$\Gamma_{Z_2} \sim 2.8 \text{ GeV}$ and v_α^k, a_α^k are those of Table 1. However, it is possible that one

(or even two) extra gauge bosons survive at low energies. In this case, M_3 (and M_4) are in general free parameters entering S_q and $A_q^{(*)}$, and as such they will be used in the present analysis. The corresponding widths $\Gamma_{3,4}$ can be calculated as functions of the masses $M_{3,4}$ and the assumed light particle content (see next section). We follow the lines of Ref. 10) for a general parametrization of v_α^k, a_α^k .

In any model with a low-energy $SU(3)_C \times SU(2)_L \times U(1) \times U(1) \times U(1)$ gauge structure, the neutral current Lagrangian can be written as

$$\mathcal{L}_{NC} = J_{\beta\mu} C_{\beta\alpha} Z_\alpha^\mu, \quad (2.7)$$

where $J_{\beta\mu} = \bar{\psi}_L^k \gamma_\mu T_\beta^k \psi_L^k$ [all fields are taken to be left-handed, $L \equiv \frac{1}{2}(1-\gamma_5)$] are four fixed (diagonal) neutral currents with properly normalized orthogonal generators, obtained by expanding the $SU(3)_C \times U(1)_Q$ -neutral E_6 Cartan subalgebra. We choose for convenience T_β to be T_{3L}, Y, Y', Y'' (their action on the E_6 fundamental representation is given in Table 2). Then, the model dependence not associated to the fermion assignments is carried by $C_{\beta\alpha}$, as is explained in detail in Appendix A. Here we simply summarize the results.

Any real matrix $C_{\beta\alpha}$ can be diagonalized by two orthogonal matrices:

$$C_{\beta\alpha} = O_{\beta\gamma} \tilde{g}_\gamma O'_{\gamma\alpha}. \quad (2.8)$$

In the simple case of no intermediate mass scales, O can be interpreted as the matrix rotating to the multiplicatively renormalized current basis, $\tilde{g}$ as the renormalized coupling constants and O' as the matrix diagonalizing the gauge boson mass matrix. $O, \tilde{g}$ and O' are model-dependent, but they are not completely arbitrary. One must impose that the weak charged currents be described by $SU(2)_L$ and that the electric charge be conserved. Besides, if the Higgs sector lies in $\underline{27}$ or $\overline{\underline{27}}$'s (as happens in superstring-inspired models), $C_{\beta\alpha}$ can be written as a function of six parameters (plus the gauge boson masses) in the case of two extra

*) Present low-energy data do not exclude new gauge bosons with masses as low as $\sim 100 \text{ GeV}$ (7), (9), (10).

gauge bosons. For one extra gauge boson ($\alpha = j = 1, 2, 3$), $C_{\beta j}$ depend (up to a sign) only on three parameters, one of which can essentially take only two values. From now on, we concentrate on the case of only one extra gauge boson. The expressions of $C_{\beta j}$ are given in Table 3. As explained in Appendix A1, the three free parameters are θ_1, θ_2 and λ . θ_3 is fixed, up to a sign, by $M_{2,3}$ and M_W . 0 is determined by θ_1 and θ_2 : a non-zero θ_1 can be generated by the presence of incomplete matter representations²²⁾, or by rescaling of the gauge couplings at the compactification scale²³⁾; θ_2 can essentially take only two values (see Appendix B). The coupling constants $\tilde{g}_2$ and $\tilde{g}_3$ are determined by λ , which is a number of order 1 and can be chosen >1 without loss of generality, whereas $0'$ is fixed by θ_3 . To parametrize the coefficients v_α^k and a_α^k , we must be able to express the currents J_β in (2.7) in terms of the physical fermions. There are in the 27 several fields with the same $SU(3)_C \times U(1)_{Q_{10}}$ charges. Then, new angles have to be introduced to define the observed fermions¹⁰⁾. When all these new mixing angles are fixed equal to zero, the known fermions are conventionally assigned inside the $[16]_{SO(10)} \subset [27]_{E_6}$, as is usually done. This means index 1 in Table 2. For its particular interplay with the parametrization of $C_{\beta j}$ in (2.8) we will, however, explicitly consider one of those angles. It corresponds to a global mixing angle β between $f_{1(2)}^c = (d_{1(2)}^c)_L, (e_{1(2)}^{1(2)})_L, \nu_{5(4)}^c)_L$ (where the upper index c in $f_{1(2)}^c$ means current eigenstate) in Table 2, see Appendix B:

$$\begin{pmatrix} f_1 \\ f_2 \end{pmatrix} = \begin{pmatrix} \cos \beta & \sin \beta \\ -\sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} f_1^c \\ f_2^c \end{pmatrix}, \quad (2.9)$$

where $f_{1(2)}$ stands for known (new) fermions^{*)}. From now on, index 1(2) will denote observed (new) fermions: $d_{1(2)}, e_{1(2)}$. As explained in Appendix B, it is not restrictive to consider only one of the two possible values of θ_2 in the case of large Higgs breaking (HB), since the other one can be reproduced by shifting

*) Obviously, the definition extends to $\nu_{5(4)}$ although they have not been observed.

$\beta \rightarrow \beta + \pi/2$ ^{*}). Then θ_2 takes only two values, depending on the mechanism lowering the rank of the gauge group from 6 to 5: $(c_2, s_2) = (1, 0)$ for non-Abelian flux breaking (FB), $(c_2, s_2) = (\sqrt{3/8}, \sqrt{5/8})$ in the case of HB.

Comparing (2.7) with (2.1) and using Table 2 and (2.9), v_α^k and a_α^k can be expressed as functions of $C_{\beta\alpha}$. The result is given in Table 4. As there are four independent currents, all v_α^k, a_α^k can be written as functions of four of them, e.g., $v_\alpha^u, v_\alpha^{d1}, a_\alpha^u$ and a_α^{d1} . Using now Table 3, we obtain, for the case of one extra gauge boson, the vector and axial couplings that are needed for phenomenological applications. They are given in Table 5.

Before giving numerical results for a set of representative cases, we have to specify the values of the different parameters for the models we consider. They are gathered in Table 6. Case (a) corresponds to FB, $(c_2, s_2) = (1, 0)$ and cases (b) and (c) to HB, $(c_2, s_2) = (\sqrt{3/8}, \sqrt{5/8})$. In the first case, physical quantities do not depend on β , whereas in cases (b) and (c) we choose $s\beta$ equal to 0 and 1 respectively, which we believe are the most natural values. λ can always be chosen larger than (or equal to) 1: as representative values for the case of gauge coupling unification at the compactification scale, we choose $\lambda = 1$ and $\lambda = 1.2$. Finally, as a sample of θ_1 values we choose $-\pi/2, -\pi/4, 0$ and $\pi/4$. Previously considered models ^{9), 11), 12)} are obtained fixing $\lambda = 1$, in which case physical quantities do not depend on θ_1 ⁺).

To calculate v_α^k and a_α^k and the cross-sections, one should also fix the value of θ_3 . We put it equal to zero in all cases. In fact, θ_3 is known to be small [$\theta_3 < 0.1$ ^{7), 10)}] and for the Z_3 decays into l^+l^- we are concerned with, it can be safely neglected. This small mixing can be important to obtain bounds on M_3^2 from a precise measurement of M_2^2 and for the decay of Z_3 into W^+W^- ¹⁵⁾. However, its contribution to the Z_3 width and the Z_3 -fermion couplings, which are the relevant quantities here, is at most a few per cent. Note that the absence of $Z^0 Z_3$ mixing

^{*}) Although a variation in θ_2 can sometimes mimic a variation in β , these two angles have a different physical meaning. Actually, β could be measured by measuring the $\bar{f}_1 f_2 Z_3$ couplings ($\bar{d}_1 d_2 Z_3, \dots$), which cannot be generated by any value of θ_2 if $\beta = 0$.

⁺) (a), (b), (c) correspond to (a), (c), (d) of Ref. 15) respectively. Case (b) there is discussed in the conclusions here.

does not mean that there are not interference effects: they result from the addition of the different amplitudes in Fig. 2 and could give important information when measured¹⁶⁾.

3. - NUMERICAL EXAMPLES

In this section, we plot the lepton differential cross-section, $(d\sigma/dM)^S$ in Eq. (2.4), and the asymmetry, $A^{\bar{P}P}(M)$ or $A^{PP}(M)$ depending on the collider in Eq. (2.5), for the E_6 superstring-inspired models of Table 6. Figures 3 and 4 correspond to the CERN collider $(a,b,\sqrt{s}) = (p,\bar{p},630 \text{ GeV})$; Figures 5 and 6 to the Tevatron, $(a,b,\sqrt{s}) = (p,\bar{p},2 \text{ TeV})$; Figs. 7 and 8 to LHC, $(a,b,\sqrt{s}) = (p,p,17 \text{ TeV})$; Figs. 9 and 10 to SSC, $(a,b,\sqrt{s}) = (p,p,40 \text{ TeV})$. In all cases we use the Duke and Owens parton distribution functions, set 1¹⁷⁾ and we assume $M_W = 81.8 \text{ GeV}$ ^{14),20)}, $M_2 = M_{Z_0}^{\text{th}} = 92.8$ and $\Gamma_2 = \Gamma_{Z_0}^{\text{th}} = 2.8 \text{ GeV}$ ($\theta_3 = 0$, see section 2 and Appendix A), and $\alpha(M_{Z_0}) = e^2/4\pi = (127.7)^{-1}$ ²⁴⁾. v_α^k and a_α^k are those of Tables 4 and 5. M_3 is fixed for illustration purposes equal to 150 GeV for the CERN collider, 300 GeV for the Tevatron, 1 TeV for LHC and 2 TeV for SSC. The Z_3 width, Γ_3 , is calculated in the two extreme cases: (i) Z_3 can decay only into the observed fermions (including the top quark); (ii) Z_3 can decay into three families of 27 fermions and the corresponding sfermions. This illustrates the possible range of variation of Γ_3 , which we calculate using the decay rates¹⁵⁾

$$\Gamma(Z_3 \rightarrow f\bar{f}) = \frac{M_3}{12\pi} (1-4\eta_f)^{1/2} C_f \left\{ (v_3^f)^2 + (a_3^f)^2 + 2[(v_3^f)^2 - 2(a_3^f)^2] \eta_f \right\}, \quad (3.1)$$

for fermions, and

$$\Gamma(Z_3 \rightarrow \tilde{f}_L \tilde{f}_R) = \frac{M_3}{12\pi} (1-4\eta_{\tilde{f}})^{3/2} \frac{1}{4} C_f (v_3^f \mp a_3^f)^2, \quad (3.2)$$

for sfermions, where $\eta_a \equiv (m_a/M_3)^2$ with m_a ($a = f, \tilde{f}$) the fermion and sfermion masses and the colour factor C_f equal to 1(3) for $SU(3)_C$ singlets (triplets). In doing numerical estimates, we neglect finite mass effects in the final state: they can be trivially included using Eqs. (3.1) and (3.2), but in case (i) they are negligible; in case (ii) they can be sizeable, but neglecting them we obtain an upper bound for the total width Γ_3 . In Table 7 we give numerical values for

the relevant quantities Γ_3/M_3 and $B(\ell^+\ell^-) \equiv \Gamma(Z_3 \rightarrow e^+e^-)/\Gamma_3 \equiv \Gamma(Z_3 \rightarrow \mu^+\mu^-)/\Gamma^3$. We consider cases (i) and (ii) above for models of type (a), (b), (c) and different values of θ_1 and λ . Γ_3/M_3 does not depend on M_3 (in the massless limit for the decay products) and is given in units of 10^3 , so that the numbers in Table 7 correspond to the width (in GeV) of a 1 TeV Z_3 .

In Figs. 11-15, we plot $(d\sigma/dMdy)_{M=M_3}^S$ and $A(y)$ in Eq. (2.6), for the same models and colliders and with the same assumptions as before. Finally, in Figs. 16-19, we give the total cross-section for the production of Z_3 decaying into e^+e^- (or $\mu^+\mu^-$). The number of events corresponds to integrated luminosities, $\int \mathcal{L} dt$, of 0.73 pb^{-1} (Sp $\bar{p}$ S) and 10 pb^{-1} (Sp $\bar{p}$ S + ACOL), Fig. 16, of 10 pb^{-1} (Tevatron), Fig. 17, and of 10^4 pb^{-1} (LHC and SSC), Figs. 18 and 19.

4. - CONCLUSIONS

In this paper, we re-examined the possibility of observing new neutral gauge bosons at hadron colliders. We concentrated on the expected best signature $Z_3 \rightarrow e^+e^-(\mu^+\mu^-)$. Results were given for the present and planned hadronic accelerators in the case of E_6 (superstring-inspired) models (see Figs. 3-19). A convenient parametrization of these models was considered (see Appendices).

We must note that in a particular class of E_6 models, the new Z_3 does not couple to $\ell^+\ell^-$ ^{10),15)}, except for its mixing with Z^0 which is small. They require a particular relation between λ and θ_1 ,

$$\lambda^2 = \frac{c_1 (3s_1 + c_1)}{s_1 (3c_1 - s_1)}, \quad (4.1)$$

which however does not seem to be related to any symmetry and therefore could arise only accidentally. It is obtained by demanding that for $\theta_3 = 0$ the Z_3 current aligns along $(1/\sqrt{10})(Y-3Y')$ in Table 2. Equation (4.1) can be satisfied in E_6 superstring-inspired models if there is rescaling of the couplings at the compactification scale²³⁾. On the other hand, if one assumes unification of the gauge couplings at the compactification scale, for the range of parameters we consider (4.1) cannot be fulfilled, but it can be approached. This explains why the e^+e^- cross-sections can be rather small in some cases [see Ref. 15)]. In the extreme case of a very small $Z_3 e^+e^-$ coupling and if, as it seems, one has to trigger at least one isolated charged lepton to avoid overwhelming backgrounds, we will have to rely on the W^+W^- signature¹⁵⁾ (if there is a non-zero $Z_3 Z^0$ mixing

and the decay is kinematically allowed) or on the production and decay of new (s)fermions (when kinematically accessible) for the detection of Z_3 at super-colliders.

Looking at the different figures, three immediate comments arise. First of all, the hadron colliders, through the identification of charged lepton pairs, could detect quite heavy neutral gauge bosons or impose strong limits on their existence (Figs. 16-19). Second, the width of these new gauge bosons is quite small (Table 7 and Figs. 3-10), and distinguishing among different models might require good resolution detectors. This point is raised and quantitatively discussed in Refs. 13) and 16). Third, the dependence on λ (and θ_1) is very significant. As we have already noted, it can make the e^+e^- signal negligible, but it can also enlarge it considerably. This can be understood looking at Tables 4 and 5.

In this paper we gave numerical results for the case of one extra gauge boson. For the case of two, one can easily extend our analysis. Using Appendix AII, Table 3 can be generalized, and together with Table 4 used to generalize the vector (v) and axial (a) couplings in Table 5. Starting from the general expression, Eq. (1.1), for the differential cross-section ($d\sigma/dM dy d\cos\theta^*$) in our plots, we have integrated out in turn some of the kinematical variables (θ^*, y, M). However, if real data were to become available, a fit to the differential cross-section in Eq. (1.1) would be advisable.

Finally, we must emphasize again that we have included neither possible experimental cuts and resolution, nor QCD corrections and backgrounds. Real life will be tougher, but more exciting.

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APPENDIX A

In this Appendix we give the details of the parametrization of $C_{\alpha\beta}$ in the E_6 low-energy neutral current Lagrangian

$$\mathcal{L}_{NC} = J_\alpha C_{\alpha\beta} Z_\beta \quad (A.1)$$

for the case of one (I) and two (II) extra gauge bosons^{*)}. The fermion assignments in the currents J_α are discussed in Appendix B.

I. ONE EXTRA GAUGE BOSON

In E_6 there are four diagonal $SU(3)_C \times U(1)_Q$ neutral generators: we choose them to be those of Table 2 and we denote by J_α ($\alpha = 1, 2, 3, 4$) the corresponding currents. If only one extra gauge boson (Z_3) survives at low energy [besides the photon (Z_1) and the observed $Z(Z_2)$], β in (A.1) runs from 1 to 3 and we call it k . Then $C_{\alpha k}$ are 12 model-dependent parameters. An equivalent set of parameters is obtained by writing

$$C_{\alpha k} = O_{\alpha j} \tilde{g}_j O'_{jk} \quad (\alpha=1,2,3,4; j,k=1,2,3) \quad , \quad (A.2)$$

where $O_{\alpha j}$ are the first three columns of a 4×4 orthogonal matrix, $\tilde{g}_j$ is a 3×3 diagonal matrix and O'_{jk} is a 3×3 orthogonal matrix, providing all together an equivalent set of $6+3+3 = 12$ parameters. In the absence of intermediate breaking scales, the matrices O , $\tilde{g}$ and O' have a definite physical meaning. $O_{\alpha j}$ defines the rotation from the E_6 current basis conventionally chosen to the one which is multiplicatively renormalized and $\tilde{g}_j$ defines the corresponding renormalized gauge coupling constants. O'_{jk} is the rotation diagonalizing the gauge boson mass matrix. In E_6 superstring-inspired models, the set of 12 free parameters reduces to three when the different physical constraints are taken into account. Actually,

^{*)} Our analysis is based on the results of Ref. 9). Here we use different conventions (the definition of the free parameters, the parametrization of the orthogonal matrices and the basis of generators) to simplify the applications. Some work is required to relate the two sets of conventions.

$$O_{\alpha j} \equiv \begin{matrix} T_{3L} \\ Y \\ Y' \\ Y'' \end{matrix} \begin{vmatrix} 1 & 0 & 0 \\ 0 & c_1 & s_1 \\ 0 & -s_1 c_2 & c_1 c_2 \\ 0 & s_1 s_2 & -c_1 s_2 \end{vmatrix} \left(-\frac{\pi}{2} \leq \theta_{1,2} < \frac{\pi}{2} \right), \quad (A.3)$$

depends on two angles $\theta_{1,2}$ ($c_{1,2} \equiv \cos \theta_{1,2} \equiv \sin \theta_{1,2}$) instead of six, because $SU(2)_L$ is a good symmetry at low energy and T_{3L} does not mix with other generators. Besides the remaining two currents must include the hypercharge Y , and a combination $\hat{Y}$ of the other two E_6 currents fixed by θ_2 , $\hat{Y} \equiv c_2 Y' - s_2 Y''$. The orthogonal generator $s_2 Y' + c_2 Y''$ is assumed to be broken at high energy and the associated current decouples from the low-energy theory. The mixing of Y with $\hat{Y}$ is given by θ_1 : a non-zero θ_1 can be generated by renormalization of gauge couplings in the presence of incomplete matter representations²²⁾, or if there is rescaling of gauge coupling at the compactification scale²³⁾. Moreover, as we argue in Appendix B, θ_2 can only take two different values, depending on the mechanism lowering the rank of E_6 : $c_2 = 1, s_2 = 0$ for non-Abelian flux breaking⁹⁾ and $c_2 = \sqrt{3/8}, s_2 = \sqrt{5/8}$ for a large Higgs v.e.v. of one of the standard model singlets in the $(27+\overline{27})^{11),12)}$. Out of the three coupling constants $\tilde{g}_j$, only one is not fixed, but it is in principle calculable and its order of magnitude known. If charged currents are described by $SU(2)_L$,

$$\tilde{g}_1^2 \equiv g_{SU(2)_L}^2 = \frac{G_F}{\sqrt{2}} 8 M_W^2 \quad (A.4)$$

is fixed by the measured M_W mass and by the Fermi constant G_F . The running positron electric charge e can be used to fix one combination of the other two couplings,

$$\frac{c_1^2}{\tilde{g}_2^2} + \frac{s_1^2}{\tilde{g}_3^2} = \frac{3}{5} \left(\frac{1}{e^2} - \frac{1}{\tilde{g}_1^2} \right). \quad (A.5)$$

We choose the remaining free parameter to be

$$\lambda \equiv \frac{\tilde{g}_2}{\tilde{g}_3}. \quad (A.6)$$

In the simple case $\theta_1 = 0$, $\tilde{g}_2$ and $\tilde{g}_3$ have a clear physical meaning: $\tilde{g}_2 = g_Y$, $\tilde{g}_3 = g_A$. The fact that the photon must couple to the electric charge $Q = T_{3L} + Y$ always allows us to decompose O' as

$$O'_{jk} \equiv \begin{pmatrix} s'_1 & c'_1 & 0 \\ c'_2 c'_1 & -c'_2 s'_1 & s'_2 \\ s'_2 c'_1 & -s'_2 s'_1 & -c'_2 \end{pmatrix} \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_3 & s_3 \\ 0 & -s_3 & c_3 \end{pmatrix}, \quad (A.7)$$

with

$$s'_1 = s_W \equiv \sin \theta_W = \frac{e}{\tilde{g}_1}, \quad c'_1 = c_W \equiv \cos \theta_W,$$

$$s'_2 = \sqrt{\frac{5}{3}} \frac{e s_1}{c_W \tilde{g}_3}, \quad c'_2 = \sqrt{\frac{5}{3}} \frac{e c_1}{c_W \tilde{g}_2}$$

and

$$c_3 \equiv \cos \theta_3, \quad s_3 \equiv \sin \theta_3, \quad -\frac{\pi}{2} \leq \theta_3 < \frac{\pi}{2}.$$

[We have chosen to parametrize O' in such a way that for $\theta_3 = 0$ the second column of $C_{\alpha k}$ corresponds to the standard Z^0 current. This greatly simplifies physical applications¹⁵⁾.] If the Higgs sector lies in fundamental E_6 representations, as happens in the models we consider, then θ_3 is known up to a sign:

$$\sin \theta_3 = \epsilon \sqrt{\frac{M_{Z^0}^2 - M_2^2}{M_3^2 - M_2^2}}, \quad \cos \theta_3 = \sqrt{\frac{M_3^2 - M_{Z^0}^2}{M_3^2 - M_2^2}} \quad \left(\begin{array}{l} \epsilon = \pm 1 \\ M_3^2 > M_2^2 \end{array} \right), \quad (A.8)$$

where $M_{Z^0}^2 = M_W^2/c_W^2$ is the standard model prediction for M_2 and $M_{2,3}$ are the neutral gauge boson masses. In summary, the 12 model-dependent numbers $C_{\alpha k}$ in (A.1), (A.2) and Table 3 can be parametrized in terms of two angles $\theta_{1,2}$ (although θ_2 can only take two different values) and a ratio of coupling constants λ . To determine completely the Lagrangian (A.1), the masses of the gauge bosons $M_{2,3}$ have to be specified [M_2 has already been measured¹⁴⁾], as well as the fermion assignments in the currents J_α (see Appendix B).

We now prove (A.8). In the gauge boson basis defined by $\theta_3 = 0$ in (A.7), the gauge boson mass matrix has the form:

$$\mathcal{M}_{kl}^2 = O_{kj}^{1T} \tilde{g}_j O_{j\alpha}^T X_{\alpha\beta} O_{\beta i} \tilde{g}_i O_{il}^1 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & A & C \\ 0 & C & B \end{pmatrix}, \quad (\text{A.9})$$

where X is the matrix of the vacuum expectation values with Higgses in 27 representations in the basis of Table 2,

$$X = \begin{pmatrix} x_1 & -\sqrt{\frac{3}{5}} x_1 & x_2 & x_3 \\ -\sqrt{\frac{3}{5}} x_1 & \frac{3}{5} x_1 & -\sqrt{\frac{3}{5}} x_2 & -\sqrt{\frac{3}{5}} x_3 \\ x_2 & -\sqrt{\frac{3}{5}} x_2 & -\frac{16}{15} x_1 + \frac{8}{15} x_2 + \frac{5}{3} x_5 & x_4 \\ x_3 & -\sqrt{\frac{3}{5}} x_3 & x_4 & x_5 \end{pmatrix},$$

$$x_1 = \frac{1}{2} (|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2 + |\langle \tilde{\nu}_3 \rangle|^2), \quad (\text{A.10})$$

$$x_2 = -\frac{1}{2\sqrt{15}} (|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2 - 4|\langle \tilde{\nu}_3 \rangle|^2),$$

$$x_3 = \frac{1}{2} (|\langle \tilde{\nu}_1 \rangle|^2 - |\langle \tilde{\nu}_2 \rangle|^2),$$

$$x_4 = -\frac{1}{2\sqrt{15}} (|\langle \tilde{\nu}_1 \rangle|^2 - |\langle \tilde{\nu}_2 \rangle|^2 + 5|\langle \tilde{\nu}_4 \rangle|^2 - 5|\langle \tilde{\nu}_5 \rangle|^2),$$

$$x_5 = \frac{1}{2} (|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2 + |\langle \tilde{\nu}_4 \rangle|^2 + |\langle \tilde{\nu}_5 \rangle|^2),$$

and

$$A = \frac{e^2}{s_w^2 c_w^2} x_1,$$

$$B = \frac{e^2}{c_w^2} \frac{1}{\lambda^2} \left\{ (\lambda^2 - 1)^2 c_1^2 s_1^2 x_1 + 2\sqrt{\frac{5}{3}} (\lambda^2 - 1) c_1 s_1 (c_1^2 + \lambda^2 s_1^2) (c_2 x_2 - s_2 x_3) \right\}$$

$$+ \frac{5}{3} (c_1^2 + \lambda^2 s_1^2)^2 \left[c_2^2 \left(-\frac{16}{15} x_1 + \frac{8}{15} x_2 + \frac{5}{3} x_3 \right) - 2 s_2 c_2 x_4 + s_2^2 x_5 \right] \Big\} , \quad (\text{A.11})$$

$$C = \frac{e^2}{s_W c_W^2} \frac{1}{\lambda} \left[(1 - \lambda^2) c_1 s_1 x_4 - \sqrt{\frac{5}{3}} (c_1^2 + \lambda^2 s_1^2) (c_2 x_2 - s_2 x_3) \right] .$$

Then θ_3 in (A.8) is found by diagonalizing the positive definite 2×2 submatrix in (A.9),

$$\begin{pmatrix} M_2^2 & 0 \\ 0 & M_3^2 \end{pmatrix} = \begin{pmatrix} c_3 & -s_3 \\ s_3 & c_3 \end{pmatrix} \begin{pmatrix} A & C \\ C & B \end{pmatrix} \begin{pmatrix} c_3 & s_3 \\ -s_3 & c_3 \end{pmatrix} , \quad (\text{A.12})$$

and noting that $A = M_W^2 / c_W^2 = M_{Z0}^2$. Knowing A , the gauge boson mass matrix $\mathcal{M}^2$ depends only on two quantities, B and C , or equivalently on M_2^2 and M_3^2 , allowing us to determine θ_3 as in Eq. (A.8), with $\epsilon = \text{sign}(C)$. B and C in (A.9) are not arbitrary, because C depends only on doublet vacuum expectation values ($\langle \tilde{v}_{1,2,3} \rangle$) and, therefore, is bounded by M_W , up to Clebsch-Gordan coefficients. Then knowing one of the masses, M_2 or M_3 , one can bound the other. A detailed discussion of the phenomenological implications of these bounds is given in Ref. 15).

Finally, we comment on the range of variation of $\theta_{1,2,3}$ and λ . It is not restrictive to assume that $-\pi/2 < \theta_{1,2,3} < \pi/2$, noting that the global sign of c_1, s_1 does not enter $C_{\alpha k}$ and that the sign of $J_{\alpha 0}^{\tilde{g}_j} \tilde{g}_{j0}^{\prime k}$ cannot be measured. On the other hand, we can assume $\lambda > 1$ in (A.6) because interchanging $\tilde{g}_2$ and $\tilde{g}_3$ amounts to redefining $\theta_1 (\theta_1 \rightarrow \theta_1 \pm \pi/2)$.

II. TWO EXTRA GAUGE BOSONS

This is a generalization of case I [we follow closely Ref. 10)]. In this case, $C_{\alpha\beta}$ in (A.1) is a real 4×4 matrix that can be decomposed as the product

$$C_{\alpha\beta} = O_{\alpha\rho} \tilde{g}_\rho O'_{\rho\beta} . \quad (A.13)$$

The same phenomenological constraints as for the case of one extra gauge boson apply, but now four light gauge fields must be considered. Then (in obvious notation)

$$O_{\alpha\rho} = \begin{pmatrix} 1 & 0 \\ 0 & R(\theta) \end{pmatrix} \quad (A.14)$$

where

$$R(\theta) = \begin{pmatrix} c_1 c_3 & s_1 c_3 & s_3 \\ -c_1 s_2 s_3 - s_1 c_2 & -s_1 s_2 s_3 + c_1 c_2 & s_2 c_3 \\ -c_1 c_2 s_3 + s_1 s_2 & -s_1 c_2 s_3 - c_1 s_2 & c_2 c_3 \end{pmatrix}, \quad -\frac{\pi}{2} \leq \theta_i < \frac{\pi}{2}, \quad (A.15)$$

is a general 3×3 rotation matrix which depends on three angles, $\theta_{1,2,3}$, that are fixed in any particular model. Two of the four coupling constants are now independent. (A.4) applies and (A.5) generalizes to

$$\sum_{\ell=2,3,4} \frac{O_{2\ell}^2}{\tilde{g}_\ell^2} = \frac{3}{5} \left(\frac{1}{e^2} - \frac{1}{\tilde{g}_1^2} \right). \quad (A.16)$$

The two arbitrary coupling constants can be estimated assuming unification and a definite matter content and symmetry breaking pattern. O' is a 4×4 orthogonal matrix whose first column is fixed by requiring that the photon couples to the electric charge. Then we can generalize (A.7) to

$$O'_{\alpha\beta} = \begin{pmatrix} s'_1 & c'_1 & 0 & 0 \\ c'_2 c'_1 & -c'_2 s'_1 & s'_2 & 0 \\ c'_3 s'_2 c'_1 & -c'_3 s'_2 s'_1 & -c'_3 c'_2 & -s'_3 \\ s'_3 s'_2 c'_1 & -s'_3 s'_2 s'_1 & -s'_3 c'_2 & c'_3 \end{pmatrix} \begin{pmatrix} 1 & 0 \\ \vdots & \vdots \\ 0 & R(\delta) \\ \vdots & \vdots \end{pmatrix}, \quad (A.17)$$

$$-\frac{\pi}{2} \leq \delta_i < \frac{\pi}{2},$$

where

$$s'_1 = s_w, \quad c'_1 = c_w, \quad c'_2 c'_1 = \sqrt{\frac{5}{3}} e \frac{O_{22}}{\tilde{g}_2}, \quad c'_3 s'_2 c'_1 = \sqrt{\frac{5}{3}} e \frac{O_{23}}{\tilde{g}_3}, \quad s'_3 s'_2 c'_1 = \sqrt{\frac{5}{3}} e \frac{O_{24}}{\tilde{g}_4}$$

and R is an arbitrary 3×3 orthogonal matrix [see (A.14)]. $\theta'_{1,2,3}$ are fixed if we assume $s'_2 > 0$. For $\delta_i = 0$, the second column of $C_{\alpha\beta}$ corresponds to the standard Z^0 direction. If the Higgs sector lies in $(27+27)$'s of E_6 (as is the case in superstring-inspired models) and the charged currents are described by $SU(2)_L$, only one combination of the three angles δ_i is an independent parameter. The other two possible combinations are functions of the first one and of the gauge boson masses $M_{2,3,4}$.

Summarizing, the 16 numbers $C_{\alpha\beta}$ in (A.1) depend on four angles and two coupling constants: $\theta_{1,2,3}$ and, say, $\delta_3, \tilde{g}_{3,4}$. The particular case of one extra gauge boson is obtained by fixing $\theta_3, \delta_{2,3} = 0$ and taking $M_4 \rightarrow \infty$ [δ_1 is called θ_3 in case (I)].

To prove that there is only one independent combination of δ_i , one can write down the gauge boson mass matrix in the basis $\delta_i = 0, \mathcal{M}^2 = 0, \tilde{g}_0^T X g_0$, where X is given in (A.10) and 0 and $0'$ in (A.14), (A.15) and (A.17). $\mathcal{M}^2$ is a function of the moduli of the five vacuum expectation values $\langle \tilde{v}_j \rangle$, but one combination of them is fixed, $x_1 = M_W^2 s_W^2 / e^2$. Then the diagonalizing matrix $R(\delta_i)$ and the three gauge boson masses $M_{2,3,4}$ are functions of four independent quantities. One can choose to express everything as a function of the three physical masses and one combination of mixing angles.

The range of variation of $\theta_{1,3}$ and $\delta_{1,3}$ in (A.15) and (A.17) (the signs of $c\theta_2$ and $c\delta_2$ are not fixed by the given parametrization and we choose them positive) is obtained by noting that the global sign of $c\theta_3$, $s\theta_3$ and $c\theta_1$ does not enter $C_{\alpha\beta}$ and that the sign of the currents $J_{\alpha}^0 \tilde{g}_{\alpha\rho}^0 J_{\rho}^0$ cannot be measured.

APPENDIX B

We now discuss the fermion assignments in the J_α currents entering the Lagrangian of Eq. (A.1). A more detailed discussion is given in Ref. 10). Out of the many angles which can be introduced to describe the mixing of all fermions with the same $SU(3)_C \times U(1)_Q$ quantum numbers in E_6 superstring-inspired models, we will consider only one. Its precise determination, as that of other poorly-known angles, requires the study of the interactions of the new fermions expected in these models. This is not our concern here and will be discussed elsewhere.

Matter fields lie in fundamental representations, 27, of E_6 , whose field content is given in Table 2. We assume, as suggested by the smallness of flavour-changing neutral currents, that the fermion assignments within the 27 are the same for all generations. Moreover, it is believed that a good approximation to the pattern of fermion masses can be obtained as the result of large $\Delta T_L = 0$ [$SU(2)_L$ invariant] masses and small $\Delta T_L = \frac{1}{2}$ [$SU(2)_L$ doublet] masses. In that case, three angles are required to describe the mixing within the 27, which rotate $d_{1,2}^c$, $(e_{1,2}^{1,2})_L$ and $\nu_{4,5L}$ respectively. We will only consider the common angle defined by the rotation of the current eigenstates $f_{1(2)}^c = (d_{1(2)L}^c, (e_{1(2)}^{1(2)})_L, \nu_{5(4)L})$ in Table 2 to the mass eigenstates $f_{1(2)}$,

$$\begin{pmatrix} f_1 \\ f_2 \end{pmatrix} = \begin{pmatrix} \cos \beta & \sin \beta \\ -\sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} f_1^c \\ f_2^c \end{pmatrix}, \quad (B.1)$$

fixing to zero the other two angles. $\beta = 0$ in (B.1) corresponds to the standard assignment of known fermions in the $[16]_{SO(10)} \subset [27]_{E_6}$.

Taking β into account, the E_6 superstring-inspired models with one extra gauge boson discussed in Appendix AI can be described by only two values of θ_2 . They correspond to the two known ways of going from E_6 to a low-energy group $SU(3)_C \times SU(2)_L \times U(1) \times U(1)$. One is non-Abelian flux breaking, which lowers the rank by one unit and implies $\theta_2 = 0$. The second one is Abelian flux breaking, which conserves the rank, followed by Higgs breaking at a large scale. For Higgs

breaking, we have only two $SU(2)_L \times U(1)_Y$ singlets at our disposal in the fundamental (27) representation of E_6 , $\tilde{\nu}_4$ and $\tilde{\nu}_5^*$). These two singlets form an $SU(2)_N$ doublet, where $SU(2)_N$ is an E_6 subgroup which commutes with the standard model subgroup $SU(3)_C \times SU(2)_L \times U(1)_Y$. Actually, E_6 contains $SU(6) \times SU(2)_N$ as a maximal subgroup, under which $27 = (\underline{15}, 1) + (\bar{6}, 2)$, whereas $SU(6)$ contains the standard $SU(5)$, under which $\underline{15} = \underline{10} + \underline{5}$ and $\bar{6} = \bar{5} + 1$. For the rank-6 group H surviving after compactification there are two alternatives: H contains or does not contain $SU(2)_N$. In the first case, one can assume without loss of generality that $\langle \tilde{\nu}_4 \rangle = 0$ and $\langle \tilde{\nu}_5 \rangle \neq 0$. This configuration can be reached from any combination with $\langle \tilde{\nu}_4 \rangle \neq 0$, $\langle \tilde{\nu}_5 \rangle \neq 0$ by means of a suitable $SU(2)_N$ gauge transformation. Therefore, using the quantum number assignments in Table 1, one finds $c_2 = \sqrt{3/8}$, $s_2 = \sqrt{5/8}$. In the second case, if one starts from $\langle \tilde{\nu}_4 \rangle \neq 0$, $\langle \tilde{\nu}_5 \rangle \neq 0$, the rotation considered before is no longer allowed, because $SU(2)_N$ is broken and $\tilde{\nu}_4$ and $\tilde{\nu}_5$ have different quantum numbers under the residual group H . Therefore, since we are interested in lowering the rank only by one unit, we have to assume that either $\tilde{\nu}_4$ or $\tilde{\nu}_5$ acquires a non-vanishing v.e.v., but not both. The two cases are equivalent, as we prove now. Using Eq. (B.1), one can see that any model with $\langle \tilde{\nu}_4 \rangle = 0$, $\langle \tilde{\nu}_5 \rangle \neq 0$ and given β is equivalent to a model with $\langle \tilde{\nu}_4 \rangle \neq 0$, $\langle \tilde{\nu}_5 \rangle = 0$ and $\beta \rightarrow \beta + \pi/2$. Therefore, it not restrictive to assume $\langle \tilde{\nu}_4 \rangle = 0$, $\langle \tilde{\nu}_5 \rangle \neq 0$, and then $c_2 = \sqrt{3/8}$, $s_2 = \sqrt{5/8}$. Note that, if one is only concerned with the standard quarks and leptons $[d_L^C, (e^V)_L] \in f_1$ in Eq. (B.1), a variation in θ_2 can mimic a variation in β by properly redefining the other parameters. However, these two angles are physically different. The difference becomes apparent only when the exotic quarks and leptons $[D_L^C, (\bar{H}^0)_L] \in f_2$ in Eq. (B.1) are considered. For example, $\beta \neq 0, \pi/2$ implies a non-zero $\overline{D}_L^C d_L^C Z_3$ coupling, which can never be reproduced varying θ_2 if $\beta = 0, \pi/2$.

Summarizing, in the models under consideration we allow only for one non-zero angle β in the fermion assignments within the 27 of E_6 . This removes the possible double degeneracy for θ_2 in the case of a large Higgs v.e.v..

*) θ_2 is arbitrary in E_6 models with arbitrary Higgs sectors.

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TABLE 1

Standard model couplings of the known fermions ($u \equiv u, c, t$; $d \equiv d, s, b$; $e \equiv e, \mu, \tau$; $\nu \equiv \nu_e, \nu_\mu, \nu_\tau$) to the electroweak gauge bosons γ and Z^0 . Here and in the following, we adopt the definitions $\gamma \equiv s_W A_3 + c_W B$ and $Z^0 \equiv c_W A_3 - s_W B$, where $s_W \equiv \sin\theta_W$, $c_W \equiv \cos\theta_W$ and A_3 and B are the gauge bosons associated to T_{3L} and Y , respectively (see Table 2). Finally, e denotes the (running) positron charge.

	γ	Z^0
ν^u	$\frac{2}{3}e$	$\frac{e}{s_W c_W} (\frac{1}{4} - \frac{2}{3} s_W^2)$
a^u	0	$-\frac{e}{s_W c_W} \frac{1}{4}$
ν^d	$-\frac{1}{3}e$	$\frac{e}{s_W c_W} (-\frac{1}{4} + \frac{1}{3} s_W^2)$
a^d	0	$\frac{e}{s_W c_W} \frac{1}{4}$
ν^e	$-e$	$\frac{e}{s_W c_W} (-\frac{1}{4} + s_W^2)$
a^e	0	$\frac{e}{s_W c_W} \frac{1}{4}$
ν^ν	0	$\frac{e}{s_W c_W} \frac{1}{4}$
a^ν	0	$-\frac{e}{s_W c_W} \frac{1}{4}$

TABLE 2

Action of the orthogonal basis ($T_{3L}, Y, Y' Y''$) for E_6 neutral currents on the fundamental representation, 27.

	T_{3L}	$\sqrt{\frac{5}{3}} Y$	$\sqrt{\frac{5}{3}} Y'$	Y''
$\begin{pmatrix} u \\ d_1 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	1/6	1/3	0
d_{2L}	0	-1/3	-2/3	0
u_L^c	0	-2/3	1/3	0
d_{1L}^c	0	1/3	-1/6	1/2
d_{2L}^c	0	1/3	-1/6	-1/2
$\begin{pmatrix} \nu_1 \\ e_1 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	-1/2	-1/6	1/2
$\begin{pmatrix} \nu_2 \\ e_2 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	-1/2	-1/6	-1/2
$\begin{pmatrix} e_2 \\ \nu_3 \end{pmatrix}_L^c$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	1/2	-2/3	0
e_{1L}^c	0	1	1/3	0
ν_{4L}	0	0	5/6	-1/2
ν_{5L}	0	0	5/6	1/2

TABLE 3

Neutral current couplings for the photon, $C_{\alpha 1}$, the observed Z, $C_{\alpha 2}$, and a new (E_6) gauge boson Z_3 , $C_{\alpha 3}$, in the current basis of Table 2. We use the notation $s_{1,2,3} \equiv \sin\theta_{1,2,3}$, $c_{1,2,3} \equiv \cos\theta_{1,2,3}$ and $\lambda = \tilde{g}_2/\tilde{g}_3$.

$$C_{11} = e$$

$$C_{12} = \frac{ec_W}{s_W} c_3$$

$$C_{13} = \frac{ec_W}{s_W} s_3$$

$$C_{21} = \sqrt{\frac{5}{3}} e$$

$$C_{22} = -\sqrt{\frac{5}{3}} \frac{e}{c_W} [c_3 s_W + s_3 s_1 c_1 (\lambda - \frac{1}{\lambda})]$$

$$C_{23} = -\sqrt{\frac{5}{3}} \frac{e}{c_W} [s_3 s_W - c_3 s_1 c_1 (\lambda - \frac{1}{\lambda})]$$

$$C_{31} = 0$$

$$C_{32} = \sqrt{\frac{5}{3}} \frac{e}{c_W} s_3 (\frac{1}{\lambda} c_1^2 + \lambda s_1^2) c_2$$

$$C_{33} = -\sqrt{\frac{5}{3}} \frac{e}{c_W} c_3 (\frac{1}{\lambda} c_1^2 + \lambda s_1^2) c_2$$

$$C_{41} = 0$$

$$C_{42} = -\sqrt{\frac{5}{3}} \frac{e}{c_W} s_3 (\frac{1}{\lambda} c_1^2 + \lambda s_1^2) s_2$$

$$C_{43} = \sqrt{\frac{5}{3}} \frac{e}{c_W} c_3 (\frac{1}{\lambda} c_1^2 + \lambda s_1^2) s_2$$

TABLE 4

Vector (v) and axial (a) couplings for the matter fields in the fundamental E_6 representation, 27, as functions of the current couplings, $C_{\beta\alpha}$.

$$v_{\alpha}^u = \frac{1}{4} C_{1\alpha} + \frac{5}{12} \sqrt{\frac{3}{5}} C_{2\alpha}$$

$$a_{\alpha}^u = -\frac{1}{4} C_{1\alpha} + \frac{1}{4} \sqrt{\frac{3}{5}} C_{2\alpha} - \frac{1}{3} \sqrt{\frac{3}{5}} C_{3\alpha}$$

$$v_{\alpha}^{d1} = -\frac{1}{4} C_{1\alpha} - \frac{1}{12} \sqrt{\frac{3}{5}} C_{2\alpha} + \frac{1}{4} \sqrt{\frac{3}{5}} C_{3\alpha} - \frac{1}{4} c_{2\beta} C_{4\alpha}$$

$$a_{\alpha}^{d1} = \frac{1}{4} C_{1\alpha} - \frac{1}{4} \sqrt{\frac{3}{5}} C_{2\alpha} - \frac{1}{12} \sqrt{\frac{3}{5}} C_{3\alpha} - \frac{1}{4} c_{2\beta} C_{4\alpha}$$

$$v_{\alpha}^{e1} = -2v_{\alpha}^u - v_{\alpha}^{d1}$$

$$a_{\alpha}^{e1} = a_{\alpha}^{d1}$$

$$v_{\alpha}^{d2} = -v_{\alpha}^u - v_{\alpha}^{d1}$$

$$a_{\alpha}^{d2} = -a_{\alpha}^u - a_{\alpha}^{d1}$$

$$v_{\alpha}^{e2} = v_{\alpha}^{d1} - v_{\alpha}^u$$

$$a_{\alpha}^{e2} = -a_{\alpha}^u - a_{\alpha}^{d1}$$

$$v_{\alpha}^{v1} = -v_{\alpha}^{d1} - \frac{1}{2}v_{\alpha}^u - \frac{1}{2}a_{\alpha}^u$$

$$v_{\alpha}^{v2} = a_{\alpha}^{d1}$$

$$v_{\alpha}^{v3} = a_{\alpha}^u$$

$$v_{\alpha}^{v4} = v_{\alpha}^{d1} + \frac{1}{2}v_{\alpha}^u - \frac{1}{2}a_{\alpha}^u$$

$$v_{\alpha}^{v5} = -a_{\alpha}^u - a_{\alpha}^{d1}$$

$$a_{\alpha}^{vi} = -v_{\alpha}^{vi}, \quad i = 1, 2, \dots, 5$$

$$v_{\alpha}^{d1}d_2 = -tg2\beta \frac{1}{4}(v_{\alpha}^u + a_{\alpha}^u + 2v_{\alpha}^{d1} + 2a_{\alpha}^{d1})$$

$$v_{\alpha}^{e1}e_2 = v_{\alpha}^{v1}v_2 = v_{\alpha}^{v4}v_5 = -v_{\alpha}^{d1}d_2$$

$$a_{\alpha}^{d1}d_2 = a_{\alpha}^{e1}e_2 = a_{\alpha}^{v1}v_2 = a_{\alpha}^{v4}v_5 = v_{\alpha}^{d1}d_2$$

TABLE 5

Vector (v) and axial (a) gauge couplings for the up (u,c,t) and down (d,s,b) quarks in (E_6) superstring-inspired models with one extra gauge boson. The remaining couplings can be trivially computed using the relations of Table 4.

$$v_1^u = \frac{2}{3} e$$

$$v_2^u = \frac{e}{s_W c_W} \left[c_3 \left(\frac{1}{4} - \frac{2}{3} s_W^2 \right) - s_3 \frac{5}{12} s_W s_1 c_1 \left(\lambda - \frac{1}{\lambda} \right) \right]$$

$$v_3^u = \frac{e}{s_W c_W} \left[s_3 \left(\frac{1}{4} - \frac{2}{3} s_W^2 \right) + c_3 \frac{5}{12} s_W s_1 c_1 \left(\lambda - \frac{1}{\lambda} \right) \right]$$

$$a_1^u = 0$$

$$a_2^u = \frac{e}{s_W c_W} \left\{ -c_3 \frac{1}{4} - s_3 s_W \frac{1}{\lambda} \left[\frac{1}{4} s_1 c_1 (\lambda^2 - 1) + \frac{1}{3} (c_1^2 + \lambda^2 s_1^2) c_2 \right] \right\}$$

$$a_3^u = \frac{e}{s_W c_W} \left\{ -s_3 \frac{1}{4} + c_3 s_W \frac{1}{\lambda} \left[\frac{1}{4} s_1 c_1 (\lambda^2 - 1) + \frac{1}{3} (c_1^2 + \lambda^2 s_1^2) c_2 \right] \right\}$$

$$v_1^d = -\frac{1}{3} e$$

$$v_2^d = \frac{e}{s_W c_W} \left\{ c_3 \left(-\frac{1}{4} + \frac{1}{3} s_W^2 \right) + s_3 \frac{1}{4} s_W \frac{1}{\lambda} \left[\frac{1}{3} s_1 c_1 (\lambda^2 - 1) + (c_1^2 + \lambda^2 s_1^2) (c_2 + \sqrt{\frac{5}{3}} s_2 c_{2\beta}) \right] \right\}$$

$$v_3^d = \frac{e}{s_W c_W} \left\{ s_3 \left(-\frac{1}{4} + \frac{1}{3} s_W^2 \right) - c_3 \frac{1}{4} s_W \frac{1}{\lambda} \left[\frac{1}{3} s_1 c_1 (\lambda^2 - 1) + (c_1^2 + \lambda^2 s_1^2) (c_2 + \sqrt{\frac{5}{3}} s_2 c_{2\beta}) \right] \right\}$$

$$a_1^d = 0$$

$$a_2^d = \frac{e}{s_W c_W} \left\{ c_3 \frac{1}{4} + s_3 \frac{1}{4} s_W \frac{1}{\lambda} \left[s_1 c_1 (\lambda^2 - 1) + (c_1^2 + \lambda^2 s_1^2) \left(-\frac{1}{3} c_2 + \sqrt{\frac{5}{3}} s_2 c_{2\beta} \right) \right] \right\}$$

$$a_3^d = \frac{e}{s_W c_W} \left\{ s_3 \frac{1}{4} - c_3 \frac{1}{4} s_W \frac{1}{\lambda} \left[s_1 c_1 (\lambda^2 - 1) + (c_1^2 + \lambda^2 s_1^2) \left(-\frac{1}{3} c_2 + \sqrt{\frac{5}{3}} s_2 c_{2\beta} \right) \right] \right\}$$

TABLE 6

Values of the parameters $(\theta_1, \theta_2, \lambda, \beta)$ corresponding to the three types of models considered in the text. For $\lambda = 1$, which implies no dependence on θ_1 , popular models are recovered (last column).

	θ_1	(c_2, s_2)	λ	s_β	model ($\lambda=1$, no θ_1 dependence)
(a)	$[-\frac{\pi}{2}, \frac{\pi}{2})$	(1, 0)	[1, 1.2]	—	Minimal ⁹⁾
(b)	$[-\frac{\pi}{2}, \frac{\pi}{2})$	$(\sqrt{\frac{3}{8}}, \sqrt{\frac{5}{8}})$	[1, 1.2]	0	Intermediate ¹¹⁾ scale Higgs breaking
(c)	$[-\frac{\pi}{2}, \frac{\pi}{2})$	$(\sqrt{\frac{3}{8}}, \sqrt{\frac{5}{8}})$	[1, 1.2]	1	Intermediate ¹²⁾ scale Higgs breaking

TABLE 7

The quantities Γ_3/M_3 and $B(e^+e^-) \equiv \Gamma(Z_3 \rightarrow e^+e^-)/\Gamma_3$ for different types of models [(a),(b),(c)] and different values of the model parameters [$\lambda = 1, 1.2$; $\theta_1 = -\pi/2, -\pi/4, 0, \pi/4$]. Cases (i) and (ii) are defined in the text.

Model	(λ, θ_1)	$10^3 \times \Gamma_3/M_3$	$B(e^+e^-)\%$	$10^3 \times \Gamma_3/M^3$	$B(e^+e^-)\%$
(a)	(1,-)	6.5	3.6	38	0.6
	(1.2, $-\pi/2$)	9.3	3.6	55	0.6
	(1.2, $-\pi/4$)	7.3	7.9	41	1.4
	(1.2, 0)	4.5	3.6	27	0.6
	(1.2, $\pi/4$)	7.3	0.7	41	0.1
(b)	(1,-)	12	5.9	38	1.8
	(1.2, $-\pi/2$)	17	5.9	55	1.8
	(1.2, $-\pi/4$)	13	8.7	41	2.7
	(1.2, 0)	8.2	5.9	27	1.8
	(1.2, $\pi/4$)	13	3.7	41	1.2
(c)	(1,-)	6.5	5.4	38	0.9
	(1.2, $-\pi/2$)	9.3	5.4	55	0.9
	(1.2, $-\pi/4$)	7.3	6.0	41	1.1
	(1.2, 0)	4.5	5.4	27	0.9
	(1.2, $\pi/4$)	7.5	6.0	41	1.1

(i)

(ii)

FIGURE CAPTIONS

- Fig. 1 : Production of e^+e^- pairs in hadron collisions. "a" always denotes a proton and "b" a proton or an antiproton, depending on the collider. $\lambda^\pm$ corresponds to $e^\pm$, $\mu^\pm$ or $\tau^\pm$.
- Fig. 2 : Lowest order graphs contributing to the process $q\bar{q} \rightarrow \lambda^+\lambda^-$.
- Fig. 3 : Lepton differential cross-section $(d\sigma/dM)^S$ as a function of the lepton invariant mass M for SppS, models (a), (b), (c) and: (i) Z_3 decaying only into the observed fermions (including the top quark); (ii) Z_3 decaying into three families of 27 fermions and s-fermions. Solid lines correspond to $\lambda = 1$, dashed lines correspond to $\lambda = 1.2$ and $\theta_1 = -\pi/2, -\pi/4, 0, \pi/4$. $M_3 = 150$ GeV is assumed.
- Fig. 4 : Lepton asymmetry $A^{PP}(M)$ as a function of the lepton invariant mass M , for SppS, $M_3 = 150$ GeV, models (a), (b), (c) and Z_3 decaying only into the observed fermions (including the top quark). The figure is essentially unchanged if Z_3 decays into three families of 27 fermions and s-fermions. Solid lines correspond to $\lambda = 1$ and dashed lines to $\lambda = 1.2$, $\theta_1 = -\pi/2, -\pi/4, 0, \pi/4$.
- Fig. 5 : The same as in Fig. 3 but for the Tevatron and $M_3 = 300$ GeV.
- Fig. 6 : The same as in Fig. 4 but for the Tevatron and $M_3 = 300$ GeV.
- Fig. 7 : The same as in Fig. 3 but for the LHC and $M_3 = 1$ TeV.
- Fig. 8 : Lepton asymmetry $A^{PP}(M)$ as in Fig. 4 but for the LHC, $M_3 = 1$ TeV and Z_3 decaying into three families of 27 fermions and s-fermions. The figure is essentially unchanged if Z_3 decays only into the observed fermions (including the top quark).
- Fig. 9 : The same as in Fig. 3 but for the SSC and $M_3 = 2$ TeV.
- Fig. 10 : The same as in Fig. 8 but for the SSC and $M_3 = 2$ TeV.
- Fig. 11 : Lepton differential cross-section at the Z_3 peak $(d\sigma/dMdy)_{M=M_3}$ as a function of the rapidity for model (a) and: (i) Z_3 decaying only into the observed fermions, including the top quark (three upper lines);

(ii) Z_3 decaying into three families of 27 fermions and s-fermions (three lower lines). Solid lines correspond to $\lambda = 1$, θ_1 arbitrary and $\lambda = 1.2$, $\theta_1 = +\pi/4, -\pi/4$. For the Sp $\bar{p}$ S, $M_3 = 150$ GeV; for the Tevatron, $M_3 = 300$ GeV; for the LHC, $M_3 = 1$ TeV; for the SSC, $M_3 = 2$ TeV. The figures corresponding to models (b) and (c) are quite similar.

Fig. 12 : Lepton asymmetry at the Z_3 peak as a function of the rapidity $A(y)$ for the Sp $\bar{p}$ S, $M_3 = 150$ GeV and models (a), (b), (c). Solid lines correspond to $\lambda = 1$, θ_1 arbitrary and $\lambda = 1.2$, $\theta_1 = 0, -\pi/2$ for: (i) Z_3 decaying only into the observed fermions, including the top quark; (ii) Z_3 decaying into three families of 27 fermions and s-fermions. For $\lambda = 1.2$, $\theta_1 = \pi/4, -\pi/4$, dashed [dotted-dashed] lines correspond to case (i) [(ii)].

Fig. 13 : The same as in Fig. 12 but for the Tevatron and $M_3 = 300$ GeV.

Fig. 14 : The same as in Fig. 12 but for the LHC and $M_3 = 1$ TeV.

Fig. 15 : The same as in Fig. 12 but for the SSC and $M_3 = 2$ TeV.

Fig. 16 : Total cross-section $\sigma(p\bar{p} \rightarrow Z_3 \rightarrow \ell^+\ell^-)$ for Sp $\bar{p}$ S and models (a), (b), (c). Solid lines correspond to $\lambda = 1$, θ_1 arbitrary and: (i) Z_3 decaying only into the observed fermions; (ii) Z_3 decaying into three families of 27 fermions and sfermions. Dashed [dotted-dashed] lines correspond to case (i) [(ii)] and $\lambda = 1.2$, $\theta_1 = -\pi/2, -\pi/4, 0, \pi/4$. The expected number of events is given for an integrated luminosity of 10 pb^{-1} corresponding to Sp $\bar{p}$ S + ACOL. The horizontal line corresponds to $\sigma(p\bar{p} \rightarrow Z_3 \rightarrow e^+e^-) = 3\text{pb}$, the present upper bound (at 90% c.l.) from UA1 and UA2 data²⁵⁾: it can be easily translated into a lower bound on M_3 for the different models.

Fig. 17 : The same as in Fig. 16 but for the Tevatron and an integrated luminosity of 10 pb^{-1} .

Fig. 18 : The same as in Fig. 16 but for the LHC and an integrated luminosity of 10^4 pb^{-1} .

Fig. 19 : The same as in Fig. 16 but for the SSC and an integrated luminosity of 10^4 pb^{-1} .

$\sqrt{s} = ab$ center of mass energy

$$M^2 = (p_{l^+} + p_{l^-})^2$$

$$\frac{M}{2}(e^y - e^{-y}) = (p_{l^+} + p_{l^-})_{||}$$

$\theta^* = a l^-$ scattering angle in the $l^+ l^-$ center of mass

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11

Fig. 12

Fig. 13

Fig. 14

Fig. 15

Fig. 16

Fig. 17

Fig. 19

Fig. 18