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ON THE MASS AND THE SIGNATURE OF A NEW Z

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ABSTRACT

A precise measurement of the Z mass at SLC and LEP can provide information on the mass of a new neutral gauge boson, guiding its detection at hadron colliders. For a significant range of parameters the decay rate of the new Z into W^+W^- is greater than the one into $e^+e^-(\mu^+\mu^-, \tau^+\tau^-)$, giving rise to a new signature. We discuss these two issues in detail for E_6 superstring-inspired models.

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1. - INTRODUCTION

The standard model of electroweak interactions is by now well established in the energy range below 100 GeV¹⁾. Moreover, the precise comparison between the measured mass of the Z gauge boson, $M_Z^{*})$, and the standard model prediction, M_{Z0} , can also shed light on possible new physics in the TeV region. In particular, if extensions of the standard model with an extra neutral gauge boson Z_3 (of mass M_3) are considered, the difference $M_{Z0} - M_Z$ can provide an indirect constraint on the value of M_3 . On the other hand, if a new gauge boson exists, knowing its signature is of crucial importance for guiding experimental searches. Even if the coupling δ of Z_3 with a pair of W bosons goes in general to zero as $(M_W/M_3)^2$ for increasing M_3 , the fact that in the decay rate $\Gamma(Z_3 \rightarrow W^+W^-)$ there is a kinematical factor which correspondingly grows as $(M_3/M_W)^4$ can make this decay channel a very important one for the detection of the $Z_3^{+)}$. Here these issues are discussed for E_6 superstring-inspired models. A detailed discussion of these models, together with other aspects of the Z_3 physics at hadron colliders, is presented in a longer companion paper²⁾.

Extensions of the standard model with extra gauge bosons have been under consideration in the literature for a long time³⁾. In particular, new neutral gauge bosons are predicted by grand unified models based on $SO(10)$ and E_6 . Fashionable superstring models, preserving $N = 1$ supersymmetry after compactification to four dimensions, may require those gauge groups and therefore extra gauge bosons⁴⁾⁻⁷⁾. On the other hand, present experimental limits on the mass of new neutral gauge bosons from low-energy neutral current experiments and from the CERN $p\bar{p}$ -collider are rather weak^{8),9)}. They provide (to one standard deviation) a typical lower bound on the mass of an extra neutral gauge boson 0(150) GeV, although lower masses cannot be excluded in a model independent way. Future e^+e^- colliders (SLC and LEP) can look for indirect effects of the new bosons on the Z_2 mass and decays¹⁰⁾, whereas present (SppS and Tevatron) and planned (LHC and SSC) hadron colliders can aim at the direct detection of a new

*) We denote with Z_i the electroweak neutral gauge bosons of definite mass M_i : $i = 1$ corresponds to the photon ($M_1 = 0$), $i = 2$ to the observed Z, $i = 3$ to a possible extra Z.

+) We thank G. Kane for calling our attention to this point.

neutral gauge boson with a mass up to a few TeV. It is then important to have indications on the mass of a possible Z_3 and on its signature at hadron colliders. We now discuss both issues in turn. Although our analysis deals mainly with E_6 superstring-inspired models, it applies to a much larger class of extended electroweak theories.

2. - ON THE MASS AND THE MIXING OF A NEW Z

Any extension of the standard model with one extra neutral gauge boson gives rise, after removing the photon, to a mass matrix $\mathcal{M}^2$ of the form²⁾

$$\mathcal{M}^2 = \begin{pmatrix} A & C \\ C & B \end{pmatrix}, \quad (1)$$

where the first entry corresponds to the standard current eigenstate $Z^0 \equiv s_w B - c_w W_3$ (B and W_3 are associated to Y and T_{3L} , respectively and $s_w \equiv \sin\theta_w$, $c_w \equiv \cos\theta_w$) and the second one to the new Z . Such a matrix is diagonalized by means of an orthogonal transformation,

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

where we assume, for definiteness, that the eigenvalues satisfy $M_3^2 > M_2^2$. The mixing angle $\theta_3 \in [-\pi/2, \pi/2)$ (the subindex 3 is introduced for later convenience) is given by

$$s_3 \equiv \sin\theta_3 = \text{sign}(C) \sqrt{\frac{A - M_2^2}{M_3^2 - M_2^2}}, \quad c_3 \equiv \cos\theta_3 = \sqrt{\frac{M_3^2 - A}{M_3^2 - M_2^2}}. \quad (3)$$

At first sight A , B and C or, equivalently, M_2^2 , M_3^2 and θ_3 seem to be three independent parameters, the only constraint on them coming from the fact that the matrix $\mathcal{M}^2$ must be positive definite. However, this is not the case in general, and in particular in models where the Higgs sector contains only standard $SU(2)_L \times U(1)_Y$ doublets and singlets, transforming non-trivially under the new gauge interactions (as it happens in E_6 superstring-inspired models). This will allow us to extract information on M_3 from measurements on M_2 . The main point is

that A and M_2^2 can be measured (B can be traded for M_2^2 if $M_2^2 \neq A$), whereas C , which is a model-dependent quantity, is bounded from above. This bound is translated into an upper bound on M_3 , for

$$M_3^2 = A + \frac{C^2}{A - M_2^2}. \quad (4)$$

Equation (4) follows from Eq. (2), noting that

$$M_2^2 + M_3^2 = A + B, \quad (5a)$$

$$M_2^2 \cdot M_3^2 = AB - C^2. \quad (5b)$$

Let us then discuss the determination of the right-hand side of Eq. (4)*).

In the basis we chose for $\mathcal{M}^2$ in (1):

$$A = M_{Z^0}^2 = \frac{M_W^2}{C_W^2} = \frac{\pi \alpha}{\sqrt{2} G_F S_W^2} \cdot \frac{1}{C_W^2}, \quad (6)$$

where G_F is the Fermi constant and α the electromagnetic fine structure constant. That is the standard model prediction for M_2^2 ^{†)}. We assume that at the scales under consideration charged currents are described by $SU(2)_L$. Therefore, A can be measured by measuring M_W ^{&)}. On the other hand, M_2^2 is expected to be measured with high precision at SLC and LEP. Thus, knowledge of C in Eq. (4) will fix M_3^2 .

*) A study of these constraints in a specific model was performed in Ref. 6).

†) Of course, the precise standard model prediction M_{Z^0} depends, via radiative corrections, on unknown parameters like the top and the Higgs masses. We assume here three fermion generations and $m_{top} \sim 40$ GeV, taking into account only the leading standard model corrections to α ¹¹⁾: $\alpha = \alpha(M_Z) = [127.7]^{-1}$. Other smaller radiative corrections to A , due to different values of m_{top} , the Higgs sector, exotic fermions, supersymmetric partners and extra gauge bosons, will be neglected in the present analysis, where we focus on tree-level mixing effects.

&) Another determination of $M_{Z^0}^2$ can come from a (model-dependent) fit of s_W^2 to low-energy neutral-current data ^{8), 9)}.

Before giving the model-dependent form of C , we introduce some notation for the superstring models we will be using as examples.

To parametrize a general E_6 superstring-inspired model with one extra neutral gauge boson^{*}), one has to introduce two angles, θ_1 and θ_2 . Given the orthogonal basis of Table 1, (T_{3L}, Y, Y', Y'') , for flavour-conserving neutral currents in E_6 , θ_2 determines the combination $\hat{Y} \equiv \cos\theta_2 Y' - \sin\theta_2 Y''$ which remains unbroken at low energy, while θ_1 parametrizes a possible mixing of the standard Y current with the new current $\hat{Y}^{+}$). One needs also three coupling constants, $\tilde{g}_1 \equiv g_{SU(2)_L}$, $\tilde{g}_2$ and $\tilde{g}_3$, associated to the multiplicatively renormalized currents: in the simplest case where $\theta_1 = 0$, one has $\tilde{g}_2 = g_Y$ and $\tilde{g}_3 = g_{\hat{Y}}$. However, only one combination of them is a new independent parameter, which we choose to be $\lambda \equiv \tilde{g}_2/\tilde{g}_3$. Finally, C only depends on the $SU(2)_L \times U(1)_Y$ doublet vacuum expectation values through a definite combination. We find²⁾

$$C = \frac{s_w}{\lambda} M_Z^2 \left\{ (1-\lambda^2) s_1 c_1 + (c_1^2 + \lambda^2 s_1^2) \left[\frac{c_2}{3} (-4+5\xi) + \sqrt{\frac{5}{3}} s_2 \xi (-1+2\eta) \right] \right\}, \quad (7)$$

where $s_{1,2} \equiv \sin\theta_{1,2}$, $c_{1,2} \equiv \cos\theta_{1,2}$ and (see Table 1 for notation)

$$\xi \equiv \frac{|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2}{|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2 + |\langle \tilde{\nu}_3 \rangle|^2}, \quad \eta \equiv \frac{|\langle \tilde{\nu}_1 \rangle|^2}{|\langle \tilde{\nu}_1 \rangle|^2 + |\langle \tilde{\nu}_2 \rangle|^2} \quad (\xi, \eta \in [0, 1]). \quad (8)$$

For conventional assignments of quarks and leptons inside the $[16]_{SO(10)} \subset [27]_{E_6}$, $\tilde{\nu}_L \equiv \tilde{\nu}_1$, $\bar{H}^0 \equiv \tilde{\nu}_2$, $H^0 \equiv \tilde{\nu}_3$, where $\bar{H}(H)$ denotes the Higgs of hypercharge $Y = -\frac{1}{2}(\frac{1}{2})$ which couples to down (up) quarks. In general, standard left-handed sneutrinos are a combination of the $\tilde{\nu}_{1,2}$ directions²⁾

$$\begin{pmatrix} \tilde{\nu}_L \\ \bar{H}^0 \end{pmatrix} = \begin{pmatrix} c\beta & s\beta \\ -s\beta & c\beta \end{pmatrix} \begin{pmatrix} \tilde{\nu}_1 \\ \tilde{\nu}_2 \end{pmatrix}, \quad (9)$$

^{*}) A detailed discussion of this parametrization can be found in Ref. 2). See also Ref. 9).

⁺) It is a usual practice to disregard θ_1 and vary θ_2 , but in superstring-inspired models, θ_2 (in contrast with θ_1) can essentially take only two values, as we argue in Ref. 2).

where $s\beta \equiv \sin\beta$, $c\beta \equiv \cos\beta$. Then lepton number conservation, $\langle \tilde{\nu}_L \rangle = 0$, implies $\eta = (s\beta)^2$. The conventional assignment corresponds to $\beta = 0$ ($\eta = 0$), but $\beta(\eta)$ can be different from zero in general. Let us finally note that Eq. (9) is part of a common rotation of the current eigenstates $f_{1(2)}^c \equiv (d_{1(2)}^c)_L, (e_{1(2)}^{v1(2)})_L, \nu_{5(4)}^{v5(4)}_L$ of Table 1

$$\begin{pmatrix} f_1 \\ f_2 \end{pmatrix} = \begin{pmatrix} c\beta & s\beta \\ -s\beta & c\beta \end{pmatrix} \begin{pmatrix} f_1^c \\ f_2^c \end{pmatrix}, \quad (10)$$

which defines²⁾ the standard (f_1) and exotic (f_2) fermions. From now on $d_{1(2)}$ and $e_{1(2)}$ will denote the observed (exotic) $Q = 1/3$ quarks and charged leptons respectively. Summarizing, any E_6 superstring-inspired model with an extra Z_3 is characterized by three angles $(\theta_1, \theta_2, \beta)$ and a ratio of coupling constants (λ). We neglect here other possible mixing angles associated to the fermion mass matrices⁹⁾.

Using Eqs. (4), (6) and (7), constraints on the parameter space (M_2, M_3, M_{Z_0}) can be extracted for every E_6 superstring-inspired model. We work out explicitly four representative models^{4), 5)}: (a) the minimal one⁶⁾; (b) a model with no Z_3 coupling to the standard charged leptons except through its mixing with $Z^{09)}$; (c) and (d) two extreme intermediate-scale-breaking models^{5), 7)}. They exemplify the typical range of variation of the parameters entering Eq. (7). The values of the different parameters for these models are given in Table 2. θ_2 can essentially take only two values, depending on the symmetry breaking mechanism which lowers the rank of the gauge group from 6 to 5: $(c_2, s_2) = (1, 0)$ for non-Abelian flux breaking (FB) and $(c_2, s_2) = (\sqrt{3/8}, \sqrt{5/8})$ for a large Higgs breaking (HB) along the $\tilde{\nu}_5$ direction in Table 1. We discuss this point in detail in Ref. 2). Models (a) and (b) are examples of FB, while models (c) and (d) are two extreme cases of HB. In the first two cases C does not depend on β , since $\theta_2 = 0$, whereas in cases (c) and (d) $s\beta = 0$ and 1 respectively. $\lambda = 1$ characterizes models (a) and (d), and we choose the representative values $\lambda = 2$ in model (b) and $\lambda = 1$ in model (c). For $\lambda = 1$, θ_1 has no physical consequences, but in case (b) θ_1 is different from zero^{*}) and fixed by $\lambda^2 = c_1(c_1 + 3s_1)/s_1(3c_1 - s_1)^2$.

^{*}) A non-zero θ_1 can be generated by renormalization of gauge couplings in presence of incomplete matter representations [or if there is rescaling of gauge couplings at the compactification scale¹²⁾].

Finally ξ , in principle, can vary between 0 and 1. However, for $\xi = 0(1)$ the down (up) quarks would be massless. Moreover, detailed model dependent calculations of radiative electroweak breaking restrict $\xi \in [0.04, 0.27]$ in model (a)⁶⁾ and $\xi \in [0.04, 0.41]$ in model (d)⁷⁾. In case (b) we favour $0 < \xi < 1/2$, as suggested by the fact that the b quark is lighter than the t quark, and C is independent of ξ in case (c).

Since direct measurements of M_2 and M_W already exist at the CERN collider, we will assume in what follows that both masses lie within the one standard deviation ellipse of Fig. 1, obtained averaging the UA1 and UA2 determinations¹³⁾:

$$\begin{aligned} M_2 &= 92.6 \pm 1.7 \text{ GeV}, \\ M_W &= 81.8 \pm 1.5 \text{ GeV}. \end{aligned} \tag{11}$$

Figure 1 also displays the line corresponding to the standard model prediction (6) for (M_{Z^0}, M_W) , varying the angle θ_W . The values of (M_2, M_W) compatible with (11) and with the existence of a Z_3 are represented by the shaded area. We have checked that these values can be obtained in any model for a suitable pattern of VEVs. In Fig. 2 we plot Eq. (4) for the four representative models in Table 2. We fix M_{Z^0} by choosing M_W equal to its central value at the CERN collider [Eq. (11)] $M_W = 81.8 \text{ GeV}$, whereas the range in M_2 corresponds to the 1σ -limit given in Fig. 1. M_2^2 has to be smaller than $M_{Z^0}^2$ as required by Eq. (5). In the limiting case in which M_2^2 coincides with the standard value M_Z^2 , no constraint on M_3^2 can be drawn. This happens in model (a) for $\xi = 0.8$, in model (b) for $\xi = 1$ and in model (d) for $\xi = 0.4$ [it is not possible in model (c), where C is independent of ξ and non-zero]. Then $C = 0$ and Z_3 and Z^0 do not mix. This is possible only in model (c) within the favoured range of ξ . But if M_2 happens to be smaller than M_{Z^0} , an upper bound on M_3 is obtained. For instance, for $M_W = 81.8 \text{ GeV}$ and our favoured range of ξ , if M_2 were 92 GeV, then our bounds would be $M_3 < 435 \text{ GeV}$ (a), $M_3 < 311 \text{ GeV}$ (b), $M_3 = 290 \text{ GeV}$ (c) and $M_3 < 264 \text{ GeV}$ (d).

3. - ON THE W^+W^- SIGNATURE OF A NEW Z

If a new gauge boson Z_3 exists, it is essential to know its signature for direct detection at hadron colliders. It can decay into pairs of fermions, $f\bar{f}$ (and possibly of sfermions $\tilde{f}\tilde{f}$ if a supersymmetric model is considered) indepen-

dently of a possible sizeable mixing with Z^0 . In addition, Z_3 can decay into W -pairs if $M_3 > 2M_W$ and it has a non-zero mixing with Z^{0*}). (In this case, a precise measurement of M_2 at future e^+e^- colliders could demand its existence, with a mass M_3 constrained as in Fig. 2). The relevant couplings of a generic Z_3 are given in Fig. 3. The corresponding decay rates are found to be

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

$$\Gamma(Z_3 \rightarrow \tilde{f}_L \tilde{f}_L^*) = \frac{M_3}{12\pi} (1-4\eta_{\tilde{f}})^{3/2} \frac{1}{4} C_f (v_3^{\tilde{f}} + a_3^{\tilde{f}})^2, \quad (13)$$

$$\Gamma(Z_3 \rightarrow W^+ W^-) = \frac{M_3}{12\pi} (1-4\eta_W)^{3/2} (1+20\eta_W+12\eta_W^2) \eta_W^{-2} \frac{\delta^2}{16}. \quad (14)$$

In the above equations v_3^f and a_3^f are the vector and axial couplings of the generic (s)fermions f ($\tilde{f}$) to the new gauge boson Z_3 , while δ is the coupling constant associated to the vertex $Z_3 W^+ W^-$. General expressions of v_3^f , a_3^f for E_6 superstring-inspired models can be found in Table 3. Finally, $\eta_a \equiv (M_a/M_3)^2$, with $a = f, \tilde{f}, W$ and C_f is a colour factor equal to 1(3.105) for $SU(3)_C$ singlets (triplets).

As can be observed by comparing (12), (13) and (14), the $W^+ W^-$ channel can be very important for a heavy enough Z_3 (small enough η_W), unless δ is negligibly small. Actually in the physically relevant limit $M_3^2 \gg M_W^2$,

*) Decays into gaugino pairs can be safely neglected for they are suppressed by a factor $(M_W/M_3)^4$ with respect to decays into W pairs.

$$\frac{\Gamma(Z_3 \rightarrow W^+W^-)}{\Gamma(Z_3 \rightarrow \sum_f f\bar{f}, \sum_f \tilde{f}\tilde{f})} \approx \frac{1}{24} \left(\frac{M_3}{M_W} \right)^4 \frac{\delta^2}{\sum_f G_f [(v_3^f)^2 + (a_3^f)^2]}, \quad (15)$$

where sfermion masses $m_{\tilde{f}}$ have been neglected with respect to M_3 . (If sfermions are not considered, there is an enhancement factor of 3/2.)

Formulae (12+15) are model independent. If $\delta = 0(1)$, $\Gamma_{W^+W^-} \sim 0(M_3^4/M_W^4)\Gamma_{ff}$, implying a very wide (heavy) Z_3 decaying only into W pairs. But this is never the case in extended electroweak models, where δ always scales as M_3^{-2} (see below)*). Still, however, the W^+W^- sample from Z_3 decays will appear to be experimentally sizeable in hadron colliders.

In general,

$$\delta = \tilde{g}_1 c_W s_3 \quad (16)$$

This follows from the observation that W^+W^- only couple to W_3 , with strength $\tilde{g}_1$. Then, it is only through its mixing with W_3 that a new Z can couple to (and decay into) W^+W^- . However, any heavy gauge boson can mix with W_3 at most proportionally to c_W , for W_3 mixes with the photon proportionally to s_W . In the standard model, there is only one massive gauge boson, Z^0 , and it saturates c_W , then coupling to W^+W^- with strength $\tilde{g}_1 c_W$. Defining Z^0 in the same way in models with one extra gauge boson, the new Z_3 will couple to W^+W^- through its mixing with Z^0 , s_3 , and proportionally to the Z^0 coupling, $\tilde{g}_1 c_W$, proving Eq. (16). Using now Eqs. (3) and (4), we can write

$$s_3 = \frac{C}{\sqrt{(M_3^2 - A)(M_3^2 - M_2^2)}}. \quad (17)$$

This proves that, if C is bounded, as it happens in any $SU(3)_C \times SU(2)_L \times U(1)_Y$ model, then s_3 scales as M_3^{-2} . In fact, M_2^2 has been measured and A and C are essentially bounded, up to Clebsch-Gordan coefficients, by M_W^2 . This follows from

*) We thank J. Ellis for clarifying this point.

the observation that any VEV contributing to A and C has $Q = 0$ and T_{3L} and $Y \neq 0$, then contributing also to M_W^2 . In this case Eq. (15) reduces to

$$\frac{\Gamma(Z_3 \rightarrow W^+ W^-)}{\Gamma(Z_3 \rightarrow \sum_f \ell_f \bar{\ell}_f, \sum_f \tilde{\ell}_f \bar{\tilde{\ell}}_f)} \approx \frac{C^2}{M_W^4} \frac{e^2 c_W^2}{24 s_W^2 \sum_f C_f [(v_3^f)^2 + (a_3^f)^2]}, \quad (18)$$

which does not scale with M_3 : C is bounded and so is Eq. (18).

To get some insight in the size of the $W^+ W^-$ signal for E_6 superstring-inspired models, let us compare $\Gamma(Z_3 \rightarrow W^+ W^-)$ with $\Gamma(Z_3 \rightarrow e^+ e^-)$, which are easily computed using Eqs. (3)-(16) and Table 3. Figures 4 and 5 give the ratio $R \equiv \Gamma(Z_3 \rightarrow W^+ W^-)/\Gamma(Z_3 \rightarrow e^+ e^-)$ for the models (a), (b), (c), (d) of Table 2 as function of M_2 (Fig. 4) and M_3 (Fig. 5), for different values of ξ . In model (b) this ratio is very large because $\Gamma(Z_3 \rightarrow e^+ e^-)$ is negligible, but $\Gamma(Z_3 \rightarrow W^+ W^-)$ is similar to the other cases, as we show below. In fact, in this model the $W^+ W^-$ signature seems to be the only accessible one for the detection of Z_3 . For values of ξ for which there is no $Z^0 Z_3$ mixing, $\delta = 0$ and so is $\Gamma(Z_3 \rightarrow W^+ W^-)$. In all cases M_2 is required to lie in the shaded area of Fig. 1. For this range of M_2 and for given (model dependent) ξ values, M_3 cannot be arbitrarily small. Then, in some cases the corresponding line in Fig. 5 starts at values of M_3 larger than the threshold $M_3 = 2M_W$. We can observe in Figs. 4 and 5 that the ratio R depends mainly on ξ , changing little for different M_2 and M_3 values as can be guessed from Eq. (18).

From Figs. 4 and 5, it is clear that the $W^+ W^-$ signal is of experimental interest, as it is comparable to the $e^+ e^-$ one. However, in general it is not the main signal, due to the penalizing factor for W pair identification with respect to lepton pairs. The detection of $W^+ W^-$ appears to rely on the chain $Z_3 \rightarrow WW \rightarrow (\ell\nu)(jj)$ where $\ell = e, \mu$ ¹⁴⁾. Taking into account the branching ratios for $W \rightarrow e\nu, jj$ and the detector efficiency, a (optimistic) "discovery limit" of 20 $Z_3 \rightarrow W^+ W^-$ events has been estimated. In Figs. 6, 7, 8 and 9 we plot the total cross-section $\sigma(AB \rightarrow Z_3 \rightarrow W^+ W^-)$ for the CERN collider, $(A, B, \sqrt{s}) = (p, \bar{p}, 630 \text{ GeV})$, Tevatron, $(A, B, \sqrt{s}) = (p, \bar{p}, 2 \text{ TeV})$, LHC, $(A, B, \sqrt{s}) = (p, p, 17 \text{ TeV})$, and SSC, $(A, B, \sqrt{s}) = (p, p, 40 \text{ TeV})$, respectively. We also give the corresponding number of events for an integrated luminosity, $\int \mathcal{L} dt$, of 0.73 pb^{-1} (Sp\(\bar{p}\)S) and 10 pb^{-1} (Sp\(\bar{p}\)S + ACOL) in Fig. 6, of 10 pb^{-1} (Tevatron) in Fig. 7 and of 10^4 pb^{-1} (LHC and SSC) in Figs. 8 and 9. In each figure, we present plots for the models (a), (b), (c), (d) in Table 2. The calculation of the $W^+ W^-$ cross-section is standard¹⁵⁾

but we use for the Z_3 production and for its decay into W^+W^- (usually not considered) the couplings in Table 3 and Fig. 3, together with Eqs. (3)÷(16). Our parton distribution functions are those of Duke and Owens, set 1¹⁶⁾. We assume that only three standard fermion generations are kinematically accessible. If Z_3 is allowed to decay into all fermions and sfermions in the 27 representations, the cross-sections would be reduced by a factor $\sim \frac{1}{2}$. Besides a factor 2 of uncertainty characterizes the parton distribution functions. As in Fig. 5, lines corresponding to $\xi = 0, 0.5$ and 1 are shown, while the shaded areas correspond to the favoured values of ξ . As before, we fix $M_W = 81.8$ GeV and we allow M_2 to vary within the shaded area of Fig. 1. Comparing the W^+W^- cross-section with the 20 event discovery limit, we observe that to detect (or exclude) a new Z_3 through its W^+W^- signal, we have to wait for LHC and SSC (Figs. 8 and 9). The CERN collider is out of question (Fig. 6) and the Tevatron is at the edge (Fig. 7).

4. - CONCLUSIONS

We have shown, generalizing the work of Ref. 6), that if a new gauge boson Z_3 exists with the properties predicted by E_6 superstring-inspired models, the accurate measurement of the mass of the known Z_2 at SLC/LEP can provide information about its mass. A disagreement between the expected Z^0 mass and the precisely measured Z_2 mass would imply an upper bound on the mass of the new gauge boson. However, we have to emphasize that the present experimental values for M_2 and M_W are in excellent agreement with the standard model predictions (see Fig. 1). On the other hand, if SLC and LEP still find agreement with the standard model, but with smaller errors^{*)}, the present lower bound on M_3 would be strengthened (except for the very special case of exactly zero mixing).

In any case, if a new neutral gauge boson exists (with $M_3 > 2M_W$) it would generally decay into W^+W^- , through its mixing with Z^0 , at a rate comparable to its conventional decay into e^+e^- . But, given the difficulty of isolating these W^+W^- events, only in LHC/SSC this new channel can produce a clearly distinguishable signal.

When this paper was under completion we learnt that in the W sample of UA1 two events were rather unusual¹⁷⁾. In fact, it was speculated from the

*) Of course, with precise measurements of the gauge boson masses one should also wonder about all sources of (non-leading) radiative corrections which have been neglected in the present work.

kinematics of these events that they could result from the decay of a new boson of mass ~ 250 GeV into W^+W^- ¹⁸⁾. It is clear from Fig. 6 that no Z_3 arising from an E_6 superstring-inspired model can give the presently needed cross-section: they are off by more than one order of magnitude. Moreover, as explained in Section 3, the $W^+W^-Z_3$ coupling is $\tilde{g}_1^c s_3$, where $\tilde{g}_1^c$ is the Z^0 coupling and s_3 the $Z^0 Z_3$ mixing. Present experimental limits^{8),9)} require that in extended electroweak models (even more general than the E_6 superstring-inspired ones) $s_3 < 0.1$ for $M_3 \sim 250$ GeV. Therefore, the cross-section required by the $Z_3 \rightarrow W^+W^-$ interpretation of the present UA1 data is one order of magnitude larger than expected in any extended gauge theory.

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TABLE 1

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

	T_{3L}	$\sqrt{5/3} Y$	$\sqrt{5/3} Y'$	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
$\begin{pmatrix} v_1 \\ e_1 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	-1/2	-1/6	1/2
$\begin{pmatrix} v_2 \\ e_2 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	-1/2	-1/6	-1/2
$\begin{pmatrix} e_2^c \\ v_3 \end{pmatrix}_L$	$\begin{pmatrix} 1/2 \\ -1/2 \end{pmatrix}$	1/2	-2/3	0
e_{1L}^c	0	1	1/3	0
v_{4L}	0	0	5/6	-1/2
v_{5L}	0	0	5/6	1/2

TABLE 2

Values of the parameters $(\theta_1, \theta_2, \lambda, \beta)$ corresponding to the four models considered in the text. The values of ξ favoured by radiative electroweak symmetry breaking are also given. An asterisk denotes a specific value chosen for illustration, whereas a bar means that there is no dependence on that parameter.

Model	(c_1, s_1)	(c_2, s_2)	λ	ξ	$s\beta$
(a): Minimal ⁶⁾	-	(1,0)	1	[0.04, 0.27]	-
(b): $BR(Z_3 \rightarrow l^+ l^-) \ll 1$ ⁹⁾	(0.99, 0.12)	(1,0)	2*	[0, 0.5]	-
(c): Intermediate scale Higgs breaking ⁵⁾	-	$(\sqrt{3/8}, \sqrt{5/8})$	1*	-	0
(d): Intermediate scale Higgs breaking ⁷⁾	-	$(\sqrt{3/8}, \sqrt{5/8})$	1	[0.04, 0.41]	1

TABLE 3

Vector (v_3^f) and axial (a_3^f) couplings of the (s)fermions of Table 1 to the new gauge boson Z_3 . Here $d_1(2)$ and $e_1(2)$ denote the observed (exotic) $Q = 1/3$ quarks and charged leptons, respectively. For $\beta \neq 0$, $\pi/2$ off-diagonal couplings are generated, which are given at the bottom of the Table. (e is equal to $\sqrt{4\pi\alpha}$.)

$$\begin{aligned}
 v_3^u &= \frac{e}{s_w c_w} \left\{ s_3 \left(\frac{1}{4} - \frac{2}{3} s_w^2 \right) + c_3 \frac{5}{12} s_w s_1 c_1 \left(\lambda - \frac{1}{\lambda} \right) \right\} \\
 a_3^u &= \frac{e}{s_w c_w} \left\{ -s_3 \frac{1}{4} + c_3 s_w \frac{1}{\lambda} \left[\frac{1}{4} s_1 c_1 (\lambda^2 - 1) + \frac{1}{3} (c_1^2 + \lambda^2 s_1^2) c_2 \right] \right\} \\
 v_3^{d1} &= \frac{e}{s_w c_w} \left\{ s_3 \left(-\frac{1}{4} + \frac{1}{3} s_w^2 \right) - c_3 \frac{1}{4} s_w \frac{1}{\lambda} \left[\frac{1}{3} s_1 c_1 (\lambda^2 - 1) + \right. \right. \\
 &\quad \left. \left. (c_1^2 + \lambda^2 s_1^2) (\sqrt{5/3} s_2 c_2 \beta + c_2) \right] \right\} \\
 a_3^{d1} &= \frac{e}{s_w c_w} \left\{ s_3 \frac{1}{4} - c_3 \frac{1}{4} s_w \frac{1}{\lambda} \left[s_1 c_1 (\lambda^2 - 1) + (c_1^2 + \lambda^2 s_1^2) (\sqrt{5/3} s_2 c_2 \beta - \frac{1}{3} c_2) \right] \right\} \\
 \\
 v_3^{e1} &= -2v_3^u - v_3^{d1} \\
 a_3^{e1} &= a_3^{d1} \\
 v_3^{d2} &= -v_3^u - v_3^{d1} \\
 a_3^{d2} &= -a_3^u - a_3^{d1} \\
 v_3^{e2} &= v_3^{d1} - v_3^u \\
 a_3^{e2} &= -a_3^u - a_3^{d1} \\
 v_3^{v1} &= -v_3^{d1} - \frac{1}{2} v_3^u - \frac{1}{2} a_3^u \\
 v_3^{v2} &= a_3^{d1} \\
 v_3^{v3} &= a_3^u \\
 v_3^{v4} &= v_3^{d1} + \frac{1}{2} v_3^u - \frac{1}{2} a_3^u \\
 v_3^{v5} &= -a_3^u - a_3^{d1} \\
 a_3^{vi} &= -v_3^{vi}, \quad i = 1, 2, \dots, 5 \\
 \\
 v_3^{d1d2} &= -\tan 2\beta \frac{1}{4} (v_3^u + a_3^u + 2v_3^{d1} + 2a_3^{d1}) \\
 v_3^{e1e2} &= v_3^{v1v2} = v_3^{v4v5} = -v_3^{d1d2} \\
 a_3^{d1d2} &= a_3^{e1e2} = a_3^{v1v2} = a_3^{v4v5} = v_3^{d1d2}
 \end{aligned}$$

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FIGURE CAPTIONS

- Fig. 1 : Allowed region of the (M_2, M_W) plane according to the present collider data. The crosses 1 and 2 correspond to the UA1 and UA2 determinations (with statistical and systematic errors added in quadrature) while the ellipse encloses the region allowed up to one standard deviation by their weighed average. The dotted line corresponds to the standard model prediction for different values of θ_w . Only the region to the left of this line is allowed by E_6 superstring-inspired models.
- Fig. 2 : Allowed region in the (M_2, M_3) plane for the models (a), (b), (c), (d) in Table 2. Lines corresponding to $\xi = 0, 0.5$ and 1 are shown. The shaded areas correspond to the range of ξ favoured by radiative electroweak breaking. M_{Z^0} is fixed by taking $M_W = 81.8$ GeV, while M_2 is varied according to Fig. 1.
- Fig. 3 : Feynman rules for the couplings of a generic gauge boson Z_3 to fermion sfermion and W pairs.
- Fig. 4 : Allowed values of the relative branching ratio $R \equiv \Gamma(Z_3 \rightarrow W^+W^-)/\Gamma(Z_3 \rightarrow e^+e^-)$, as a function of M_2 , for the models (a), (b), (c), (d) in Table 2. The choice of parameters is the same as in Fig. 2.
- Fig. 5 : The same as Fig. 4 but for R as a function of M_3 .
- Fig. 6 : The total cross-section $\sigma(p\bar{p} \rightarrow Z_3 \rightarrow W^+W^-)$ (in nanobarns) at $\sqrt{s} = 630$ GeV, for the models (a), (b), (c), (d) in Table 2. The expected number of events is given for integrated luminosities 730 nb^{-1} (Sp̄pS) and 10 pb^{-1} (Sp̄pS + ACOL).
- Fig. 7 : The same as Fig. 6 for $\sqrt{s} = 2$ TeV and integrated luminosity 10 pb^{-1} (Tevatron).
- Fig. 8 : The same as Fig. 6 for $\sigma(pp \rightarrow Z_3 \rightarrow W^+W^-)$, $\sqrt{s} = 17$ TeV and integrated luminosity 10^4 pb^{-1} (LHC).
- Fig. 9 : The same as Fig. 8 for $\sqrt{s} = 40$ TeV and integrated luminosity 10^4 pb^{-1} (SSC).

Fig. 1

Fig. 2a

Fig. 2b

Fig. 2c

Fig. 2d

$$i \gamma^\mu (v_f^3 + a_f^3 \gamma^5)$$

$$i (K_1 - K_2)_\mu (v_f^3 + a_f^3)$$

$$i \delta [K^\tau g^{\rho\sigma} - K^\sigma g^{\rho\tau} + K_1^\rho g^{\tau\sigma} - K_1^\tau g^{\rho\sigma} + K_2^\sigma g^{\rho\tau} - K_2^\rho g^{\tau\sigma}]$$

Fig. 3

Fig. 4a

Fig. 4b

Fig. 4c

Fig. 4d

Fig. 5a

Fig. 5b

Fig. 5c

Fig. 5d

Fig. 6a

Fig. 6b

Fig. 6c

Fig. 6d

Fig. 7a

Fig. 7b

Fig. 7c

Fig. 7d

Fig. 8a

Fig. 8b

Fig. 8c

Fig. 8d

Fig. 9a

Fig. 9b

Fig. 9c

Fig. 9d