II.6 PRODUCTION OF $W_R$ AND $W_I$ BOSONS FROM SUPERSTRING-INSPIRED $E_6$ MODELS

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We study now the production of two gauge bosons, $W_R$ and $W_I$, associated with the low-energy groups arising from the breaking of $E_6$ superstring. The first model is the so-called alternative left-right model (ALRM), where the normal bounds on the $W_R$ mass do not apply, as in the standard $SU(2)_L \times SU(2)_R \times U(1)_{L-R}$ model, because of the absence of mixing with the usual $W_L$. In the second model the additional $SU(2)_L$ subgroup of $E_6$ which is obtained at relatively low energy has generators that commute with the electric charge [1]. The corresponding flavour-changing non hermitian gauge bosons $W_I$, $W_I^\dagger$ couple the conventional fermions to their exotic partners of the 27 representation of $E_6$. For both models we only focus on the flavour changing bosons $W_R$ and $W_I$ production and our results, reported in greater detail in ref. [2], show the similarity of the expected effects from the two classes of models, with discovery limits up to masses of $1.2 - 2.5$ TeV, for both SSC and LHC colliders. We also compare with previous studies of the production cross sections and, in particular, we get larger results than those obtained in ref. [3].

Within the context of all possible left-right symmetric realizations of the $E_6$ superstring, the quantum numbers of the ALRM are uniquely determined from assigning the usual fermions $(\nu_e)_L, e_L, d_L^c$ to that part of the 27 representation which transforms as a 10 under $SO(10)$ and a $\bar{5}$ under $SU(5)$, whereas the exotic counterpart of these fields, i.e. the heavy fermions $(\nu_L)_L, E_L, h^c_L$, are assigned to the $(16, \bar{5})$ term in the decomposition of the 27 of $E_6$ into $SO(10)$ and $SU(5)$ subgroups. Then, in this model, the $W_R$ has negative R-parity and nonvanishing lepton number. This means that there cannot be any mixing of the $W_R$ with the usual $W_L$. The $W_R$ boson does not couple to the $d_L^c$ quark nor the $\nu^c$ field. Hence, the usual arguments from low-energy phenomena do not constrain the mass of the charged $W_R$ boson of the ALRM model. The $W_R$ is coupled instead to the $h^c_L$ leptoquark and the $n$ field, in addition to the usual $u^c_L$ and $e^c$ particles. The coupling of $W_R$ to fermions reads

$$\mathcal{L} = \frac{1}{\sqrt{2}} g_R W_R^\mu \left( h^c \gamma_\mu u^c_L + \bar{E}^c \gamma_\mu \nu_L + \bar{\nu}^c \gamma_\mu n_L + N^c \gamma_\mu e_L \right) + h.c.$$ 

where $g_R = g_L = g$, $g$ denoting the usual $SU(2)_L$ coupling constant. The exotic fermions $h, E, N, E$ and the boson $W_R$ obtain masses from the same scale. These particles are heavy compared to the $n$ mass, which is expected to be of a few GeV order. This fact has important consequences for the $W_R$ decay modes. The dominant $W_R$ production mechanisms are $g + u \rightarrow W_R^+ + h$ and $g + \bar{u} \rightarrow W_R^- + h$. Note that the quantum numbers of $W_R$ and the conservation of R-parity imply that the production of $W_R$ from $ud$ scattering in hadronic collisions cannot take place. The production of $W_R$-pairs via the decay of a $Z'$ is forbidden as well, owing to kinematical reasons, i.e. $2 M_{W_R} > M_{Z'}$. Finally, production of the $W_R$ boson via $uh$ scattering is suppressed, owing to the smallness of the $h, u, h$ sea.

Our results are reported in Fig. 1, where we have used the distribution functions from ref. [4] with $\Lambda_{QCD} = 160$ MeV. The parton densities of Duke and Owens [5] with $\Lambda_{QCD} = 200$ MeV, lead to results differing by less than 10% from those plotted in Fig. 1. Note that with a typical branching ratio (BR) of about 1%, obtained by estimating the individual BR for the $h$ and $W_R$ particles into an observable final state to be of order 10%, we can give discovery limits for the $W_R$ mass. Assuming the minimum value for the observed cross-section at LHC to be $\sigma_{obs} \sim 10^{-4}$pb, then the $W_R^+$ could be detected up to a mass of about 2 - 2.5 TeV, and the $W_R^-$ would be observable in the range below $M_{W_R^-} = 1 - 1.5$ TeV (see Fig. 1), depending on the leptoquark mass. The $W_R$ discovery limits at SSC with a luminosity of order $10^{33}$ cm$^{-2}$ sec$^{-1}$ are not dramatically higher than those given above for LHC with luminosity of about $10^{34}$ cm$^{-2}$ sec$^{-1}$. This is summarized in Table 1. We substantially agree with the results of ref. [6].

Next, we turn our attention to the possible final state signatures. The decay modes of the leptoquark $h$ depend on the superpotential. If the $N^c_e$ is given negative R-parity, then the possible final state signatures are [6]: $jet + l^+ l^- + \bar{\nu}_T , jet + e^- + \bar{\nu}_T$, obtained from the decay modes: $h \rightarrow d + \bar{\nu}, h \rightarrow u + \bar{\nu}$, which dominate, in the assumption that sleptons are much lighter than squarks. If one assigns positive
R-parity to the \( N_e^c \), then the decay \( h \to d + \bar{N}_e \) is also possible. The \( W_R \) decay modes depend on the mass of the \( n \). This is expected to be smaller than the mass scale of the \( W_R, h, E \) and \( N^c_E \) by at least one order of magnitude [6,8]. Hence, the largely dominant decay mode is expected to be \( W_R^+ \to e^+ + n^+ R \), with an estimated branching ratio for this mode larger than 10%. This yields the possible final state signatures: \( W_R^+ \to e^+ + \gamma + p_T(\bar{\eta}) \), if the LSP is, for example, the photino \( \bar{\eta} \) and the \( n \) decays before leaving the detector, or \( W_R^+ \to e^+ + p_T(n) \), if \( n \) is a mass eigenstate and either it is the LSP and hence it is stable, owing to R-parity conservation, or it has a mean-life long enough to escape the detector before it decays. Clearly the mixing of \( n \) with \( \bar{\eta} \) and the remaining neutralinos is needed. Combining the \( h \)-decay and the \( W_R \)-decay gives rise to a final state with a very large invariant mass. Taking \( M_{W_R} = 2 \text{ TeV} \), together with the experimental value \( M_{W_L} = 80 \text{ GeV} \), one gets \( \Gamma_{tot}(W_R) = 62 (144) \text{ GeV} \) for \( n_2 = 0 \) (3).

We consider now the flavor-changing gauge boson \( W_I \) arising from the \( SU(2)_L \times U(1)_Y \times SU(2)_I \) model [3]. The dominant \( W_I \) production mechanisms are \( g + d \to W_I + h \) and \( g + d \to W_I^\dagger + \bar{h} \). As it is easier to find a \( d \)-quark, rather than a \( \bar{d} \)-quark in the proton, we expect for the production cross-sections \( \sigma(W_I) > \sigma(W_I^\dagger) \). This is clearly seen in our explicit numerical results which also confirm the similarity of the results for \( W_I \) production with the cross-section for \( W_R \) (see Fig.2). The signature of the final state is obtained combining the decay of the leptoquark with the decay of the \( W_I \) boson, as in the previous case of \( W_R \) production. The decay of \( h \) proceeds in the same way, whereas the decay of \( W_I \) yields several charged leptons, in addition to missing \( p_T \) originating from photinos and neutrinos, in the final state. Using an estimated \( BR \) of 1%, we get the discovery limits for the flavor-changing gauge bosons summarized in Table 1. We find that, at both hadron colliders, \( W_I \) will be observable up to a mass of 1.5-2 \( \text{ TeV} \), whereas for \( W_I^\dagger \) the discovery limit is given by a mass of about 1-1.5 \( \text{ TeV} \), depending on the leptoquark mass.

<table>
<thead>
<tr>
<th>( W^+ )</th>
<th>( W^- )</th>
<th>( W_I )</th>
<th>( W_I^\dagger )</th>
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<tbody>
<tr>
<td>LHC</td>
<td>2 - 2.5 TeV</td>
<td>1 - 1.5 TeV</td>
<td>1.5 - 2.2 TeV</td>
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<tr>
<td>SSC</td>
<td>2.5 - 3 TeV</td>
<td>1.2 - 2 TeV</td>
<td>2 - 2.5 TeV</td>
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Table 1

References
Fig. 1 - Cross-sections at LHC for $W^+_R$ (full) and $W^-_R$ (dashed) production as a function of $M_{W_R}$. The three different curves, from the upper to the lower curve, correspond to $m_h = 0.3, 0.6, 0.9 TeV$, respectively.

Fig. 2 - Same as Fig. 1, except for $W_I$ (dashed) and $W^I_I$ (full) production.
APPENDIX: Limits on the $W_R$ mass and on the mixing angle

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The most stringent limit on the $W_R$ mass arises from the $K_L - K_S$ mass difference $\Delta m_K$, where, in addition to the usual box diagram describing the short distance contribution to $\Delta m_K$, one gets an additional term with exchanges of a $W_L$ and a $W_R$. By assuming that the matrix element of the left-right effective Hamiltonian does not exceed the Standard Model prediction one obtains

$$8 \beta \left| (1 + \ln \frac{m_c^2}{M_L^2}) \frac{c_1 M_{LR}^1 + c_2 M_{LR}^2}{\eta M_{LL}} \right| \leq 1$$

(1)

where $\beta = (M_L/M_R)^2$, $m_c$ is the charm mass, $\eta$, $c_1$, $c_2$ are QCD coefficients ($\eta \simeq 1$, $c_1 \simeq 1.2$, $c_2 \simeq 0.08$) and

$$M_{LR}^1 = \langle \bar{K}^0 | O_{LR}^1 | K^0 \rangle = \langle \bar{K}^0 | (1 - \gamma_5) d \bar{s} (1 + \gamma_5) d | K^0 \rangle$$

(2)

$$M_{LR}^2 = \langle \bar{K}^0 | O_{LR}^2 | K^0 \rangle = \langle \bar{K}^0 | \bar{s} \gamma^\mu (1 - \gamma_5) d \bar{s} \gamma_\mu (1 + \gamma_5) d | K^0 \rangle$$

(3)

$$M_{LL} = \langle \bar{K}^0 | O_{LL} | K^0 \rangle = \langle \bar{K}^0 | \bar{s} \gamma^\mu (1 - \gamma_5) d \bar{s} \gamma_\mu (1 - \gamma_5) d | K^0 \rangle$$

(4)

A recent evaluation of these matrix elements [1] by using three-point function QCD Sum Rules gives the results that are reported in Table I together with the vacuum saturation approximation (VSA) results.

<table>
<thead>
<tr>
<th>VSA</th>
<th>QCD Sum Rules</th>
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<tbody>
<tr>
<td>$M_{LR}^1$ $(GeV^4)$</td>
<td>$10.0 \times 10^{-2}$ $(2.0 \pm 0.7) \times 10^{-2}$</td>
</tr>
<tr>
<td>$M_{LR}^2$ $(GeV^4)$</td>
<td>$-8.0 \times 10^{-2}$ $-2.0 \pm 0.7 \times 10^{-2}$</td>
</tr>
<tr>
<td>$M_{LL}$ $(GeV^4)$</td>
<td>$1.7 \times 10^{-2}$ $(1.7 \pm 0.7) \times 10^{-2}$</td>
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</tbody>
</table>

The result of this analysis is that, whereas in the case of the left-left operators the vacuum saturation approximation seems to work well, no enhancement of the left-right operators as predicted by the factorization (VSA) hypothesis is found. In other terms the three matrix elements $M_{LR}^1, M_{LR}^2, M_{LL}$ have comparable sizes when computed by QCD sum rules.

The limit for $W_R$ mass obtained using Eq.(1) is:

$$M_R \geq 700 \text{ GeV}$$

(5)

This limit holds for the economic model of pseudo-manifest left-right symmetry, where the Kobayashi-Maskawa matrices for left and right handed sectors: $U_L$ and
$U_R$ are related: $U_L = U_L K^*$, where $K$ is a diagonal unitary matrix. In more general models one can obtain less stringent limits, but only for a limited range of parameters $[3]$.  

There are a few other processes that give limits on $M_R$: $B^0 - \bar{B}^0$ mixing, $b$ decays, neutrinoless double $\beta$ decays; however, the bounds obtained by these processes are less stringent than Eq.(5).  

$W_L$ and $W_R$ can mix together by an angle $\zeta$ to form mass eigenstates $W_1$ and $W_2$:

$$W_L^\pm = \cos \zeta \ W_1^\pm - \sin \zeta \ W_2^\pm$$
$$W_R^\pm = e^{i\omega} (\sin \zeta \ W_1^\pm + \cos \zeta \ W_2^\pm).$$

There are several constraints on $\zeta$. From the theoretical bound $[4]$

$$|\zeta| \leq (M_1/M_2)^2$$

and from Eq.(5) one obtains $|\zeta| \leq 10^{-2}$. Limits of the same size come from weak universality $[5]$ and from $K_{\pi3}$ decay $[6]$, provided the CP-violating phases in $U_R$ are large. On the other hand, for small phases, one gets for $|\zeta|$ an upper limit of $2.5 \times 10^{-3}$ $[3]$.

REFERENCES