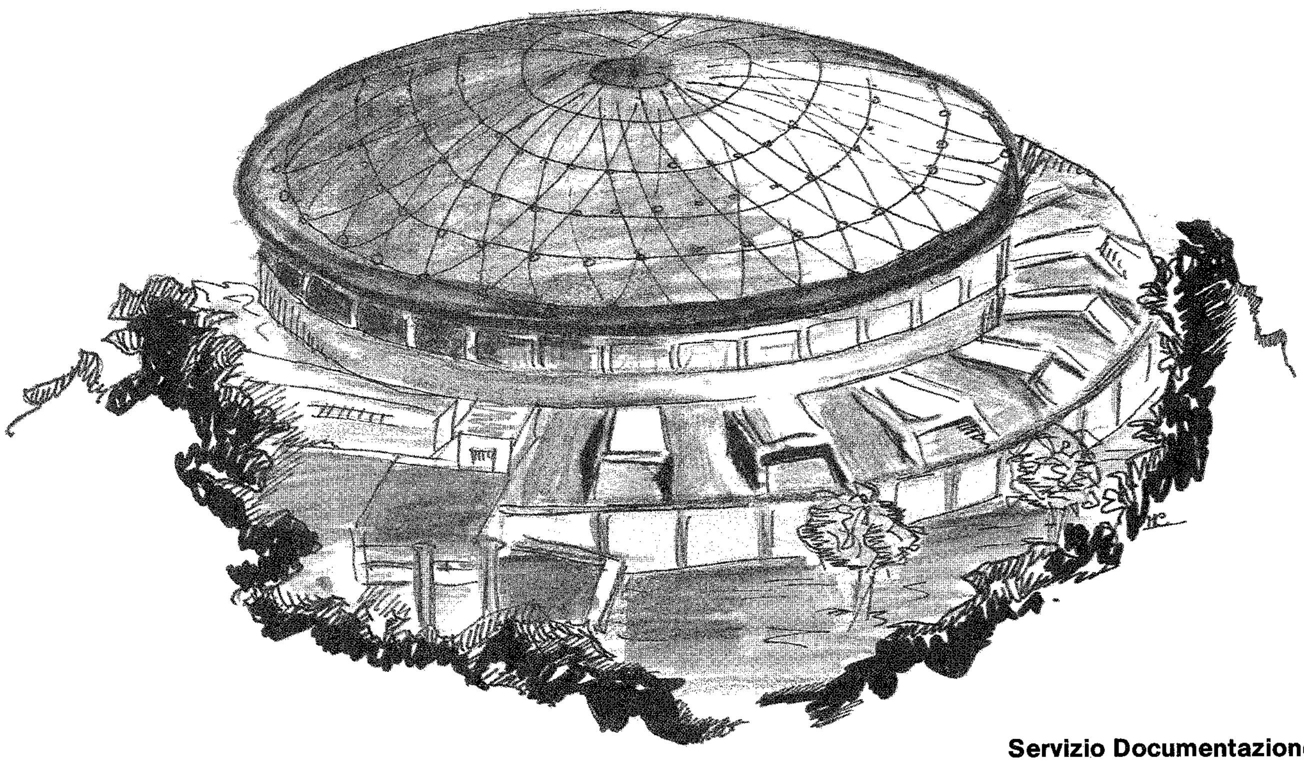
Laboratori Nazionali di Frascati

Submitted to Phys. Lett. B

LNF-90/081(PT)
29 Ottobre 1990

F. Aversa, S. Bellucci, M. Greco, P. Chiappetta:

PRODUCTION OF W_R AND W_I BOSONS FROM SUPERSTRING-INSPIRED E_6 MODELS AT HADRON COLLIDERS



Servizio Documentazione
dei Laboratori Nazionali di Frascati
P.O. Box, 13 - 00044 Frascati (Italy)

INFN - Laboratori Nazionali di Frascati
Servizio Documentazione

LNF-90/081(PT)
29 Ottobre 1990

Production of W_R and W_I Bosons from Superstring-Inspired E_6 Models at Hadron Colliders

by

Fabrizio Aversa, Stefano Bellucci, Mario Greco
INFN-Laboratori Nazionali di Frascati
P.O. Box 13
I-00044 Frascati, Italy

Pierre Chiappetta
Centre de Physique Théorique
Luminy, F-13288 Marseille, France

Abstract

We study the production cross-sections at future hadron colliders for W_R and W_I gauge bosons associated with two low-energy groups arising from the breaking of E_6 superstring. Discovery limits are given at LHC and SSC using estimated machine luminosities. Our results are compared with previous studies.

1. - Introduction and conclusions

The possibility of observing new neutral gauge bosons has been recently considered with renewed interest in the perspective of the future hadron colliders LHC and SSC. The phenomenological implications of various classes of models that provide interesting scenarios of new physics beyond the standard $SU(3) \times SU(2) \times U(1)$ have been examined in detail, particularly in the framework of superunified gauge theories.

The aim of the present paper is to study the production at LHC and SSC energies of the gauge bosons associated to two classes of superstring-inspired E_6 models, and compare the corresponding discovery limits at the expected machine luminosities. 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. This model takes advantage of the ambiguity existing in the assignment of fermions in the 27 representation of E_6 and leads to interesting phenomenological consequences at the future multi-TeV colliders, particularly if the exotic matter sector is not very heavy. In the second model the additional $SU(2)_I$ 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 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 hadron colliders. We also compare with previous studies of the production cross sections at SSC and LHC, and in particular we get larger results than those obtained in ref. [2]

2. - The left-right symmetric model

We start with the $SU(2)_L \times SU(2)_R \times U(1)_V$ model. The quantum number assignments for the 27 representations of E_6 appropriate to this model [3] are given in Table 1 from ref. [4]. Within the context of all possible left-right (LR) symmetric realizations of the E_6 superstring, the quantum numbers of the

Table 1: The quantum numbers of the 27 fermions, as given in ref. [3].

	T_{3L}	Y_L	T_{3R}	Y_R	R	B	L
N_E^c	$-\frac{1}{2}$	$-\frac{1}{6}$	$\frac{1}{2}$	$\frac{1}{6}$	$-$	0	0
E^c	$\cdot\frac{1}{2}$	$-\frac{1}{6}$	$\frac{1}{2}$	$\frac{1}{6}$	$-$	0	0
e^c	0	$\frac{1}{3}$	$\frac{1}{2}$	$\frac{1}{6}$	$+$	0	-1
e	$-\frac{1}{2}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$+$	0	1
ν_e	$\frac{1}{2}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$+$	0	1
n	0	$\frac{1}{3}$	$-\frac{1}{2}$	$\frac{1}{6}$	$-$	0	0
E	$-\frac{1}{2}$	$-\frac{1}{6}$	0	$-\frac{1}{3}$	$-$	0	0
ν_E	$\frac{1}{2}$	$-\frac{1}{6}$	0	$-\frac{1}{3}$	$-$	0	0
N_e^c	0	$\frac{1}{3}$	0	$-\frac{1}{3}$	$+(-)$	0	$-1(0)$
u	$\frac{1}{2}$	$\frac{1}{6}$	0	0	$+$	$\frac{1}{3}$	0
d	$-\frac{1}{2}$	$\frac{1}{6}$	0	0	$+$	$\frac{1}{3}$	0
h	0	$-\frac{1}{3}$	0	0	$-$	$\frac{1}{3}$	1
u^c	0	0	$-\frac{1}{2}$	$-\frac{1}{6}$	$+$	$-\frac{1}{3}$	0
h^c	0	0	$\frac{1}{2}$	$-\frac{1}{6}$	$-$	$-\frac{1}{3}$	-1
d^c	0	0	0	$\frac{1}{3}$	$+$	$-\frac{1}{3}$	0

ALRM are uniquely determined from assigning the usual fermions $(\nu_e)_L$, e_L, ${d^c}_L$ 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 $(\nu_E)_L$, 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.

The choice above interchanges the role of the fermions, with respect to the assignment leading to the conventional *LR* symmetric models [5]. This produces some distinctive physical consequences that separate this realization of the *LR* model from the conventional one. The most important inequivalence between the two realizations is to be found in the physical properties of the W_R charged boson, especially the absence of mixing between W_R and W_L. For this reason we focus, in the present work, on the production and possible detection of W_R at *LHC* and *SSC* colliders.

Strong bounds on the mass of the W_R boson in the conventional *LR* model were obtained from existing experimental data on the $K^0 - \bar{K}^0$ in ref. [6]. Taking into account short distance *QCD* corrections yields an even stronger constraint on the W_R mass [7]. However, a recent evaluation of the hadronic matrix elements using QCD sum rules [8] gives a considerably lower bound than that of refs. [6,7]. Other mass limits on W_R of the conventional LR model are coming from polarized μ-decay [9] and from other low-energy phenomena [10], but they are all less stringent than the constraint put by the $K_L - K_S$ mass difference.

In the alternative LR symmetric model ($ALRM$) 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 ${d^c}_L$ quark nor the ν^c field. Hence, the above arguments from low-energy phenomena do not constrain the mass of the charged W_R boson of the $ALRM$ model. In this model 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

$$\mathcal{L} \;=\; \frac{1}{\sqrt{2}} g_R {W_R}^{\mu} \Big(\bar{h^c}\gamma_\mu {u^c}_L \;+\; \bar{E^c}\gamma_\mu \nu_L \;+\; \bar{e^c}\gamma_\mu n_L \;+\; \bar{N^c}_E \gamma_\mu e_L \Big) \;+\; h.c. \,. \quad (1)$$

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.

3. - Associated production of W_R and leptoquark

The dominant W_R production mechanisms are $g + u \rightarrow h + {W_R}^+$ and $g + \bar{u} \rightarrow \bar{h} + {W_R}^-$. Note that the quantum numbers of W_R and the conservation of R-parity imply that the production of W_R from $u\bar{d}$ 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. $2M_{W_R} > M_{Z'}$. Finally, production of the W_R boson via $u\bar{h}$ scattering is suppressed, owing to the smallness of the $h, \bar{h}$ sea.

The differential cross-section of the process $g + u \to h + W_R^+$ reads

$$\frac{d\hat{\sigma}}{dt} = \frac{1}{16\pi s^2} < |\mathcal{M}_{LR}|^2 > , \tag{2}$$

where the amplitude is obtained from:

$$\begin{aligned} < |\mathcal{M}_{LR}|^2 > = \frac{G_F M_W^2}{3\sqrt{2}} 4\pi\alpha_s \Bigg[& - \left(\frac{t'}{s} + \frac{s}{t'} \right)\left(2 + \frac{m_h^2}{M_{W_R}^2} \right) - 2\frac{m_h^2}{M_{W_R}^2} \\ & + 2\left(2M_{W_R}^2 - m_h^2 - \frac{m_h^4}{M_{W_R}^2} \right)\left(\frac{1}{s} + \frac{1}{t'} \right) \\ & + \frac{2}{st'}\left(- \frac{m_h^6}{M_{W_R}^2} + 3m_h^2 M_{W_R}^2 - 2M_{W_R}^4 \right) \\ & + 2\frac{m_h^2}{t'^2}\left(- \frac{m_h^4}{M_{W_R}^2} - m_h^2 + 2M_{W_R}^2 \right) \Bigg] . \end{aligned} \tag{3}$$

Here we define $t' = t - m_h^2$. Next, we discuss the kinematical aspects of our calculation and give some details on the phase space.

The partonic cross-section is obtained by integrating eq. (2) in t' between the values t_1, t_2 [2]

$$\begin{aligned} t_{1,2} = & - \frac{1}{2}\left(s + m_h^2 - M_{W_R}^2 \right) \\ & \pm \frac{1}{2}\left[\left(s - m_h^2 - M_{W_R}^2 \right)^2 - 4m_h^2 M_{W_R}^2 \right]^{\frac{1}{2}} . \end{aligned} \tag{4}$$

Then, the total hadronic cross-section is deduced, as usual, by convoluting the partonic cross-section with the parton distribution functions, as

$$\sigma \sim \int_0^1 dx_1\, dx_2 \left[u^p(x_1) g^p(x_2) + g^p(x_1) u^p(x_2) \right] \hat{\sigma}(x_1 x_2 S) , \tag{5}$$

where S is the c.m. energy squared, $s = x_1 x_2 S$, and we denote by u^p, g^p the parton distribution functions relative to the proton.

In carrying out the actual computation, we use the equation

$$\sigma = \int_{-Y}^{Y} dy \int_{x_{min}}^{e^{-2|y|}} dx \left[u^p(x_1) g^p(x_2) + g^p(x_1) u^p(x_2) \right] \hat{\sigma}(x_1 x_2 S) , \qquad (6)$$

where we apply the rapidity cut given by

$$Y = min\left(2.5 , -\frac{1}{2} \ln x_{min} \right) , \qquad (7)$$

with

$$x_{min} = \frac{1}{S} \left(m_h + M_{W_R} \right)^2 . \qquad (8)$$

Furthermore, one has

$$x_{1,2} = \sqrt{x}\, e^{\pm y} . \qquad (9)$$

Note that u^p receives contributions from both valence and sea, whereas only the sea contributes to $\bar{u}^p$, so that $u^p > \bar{u}^p$. Hence, the cross-section for the associated production of ${W_R}^+$ and h is larger than the production cross-section for ${W_R}^-$ and $\bar{h}$. This is well illustrated in Fig. 1,2 for LHC and SSC energies, respectively.

For the numerical results presented here, we have used the distribution functions from ref. [11], with $\Lambda_{QCD} = 160\ MeV$. The parton densities of Duke and Owens [12] with $\Lambda_{QCD} = 200\ MeV$, lead to results differing by less than 10% from those plotted in Figs. 1 and 2. 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. With a luminosity lower by a factor 10 at SSC, with respect to LHC, one should be able to measure cross-sections with values as small as $10^{-3} pb$ at SSC. Using Fig. 2 one realizes that the W_R discovery limits at SSC with a luminosity of order 10^{33} are not dramatically higher than those given above for LHC with luminosity of about 10^{34}. This is summarized in Table 2. We substantially agree with the results of ref. [4].

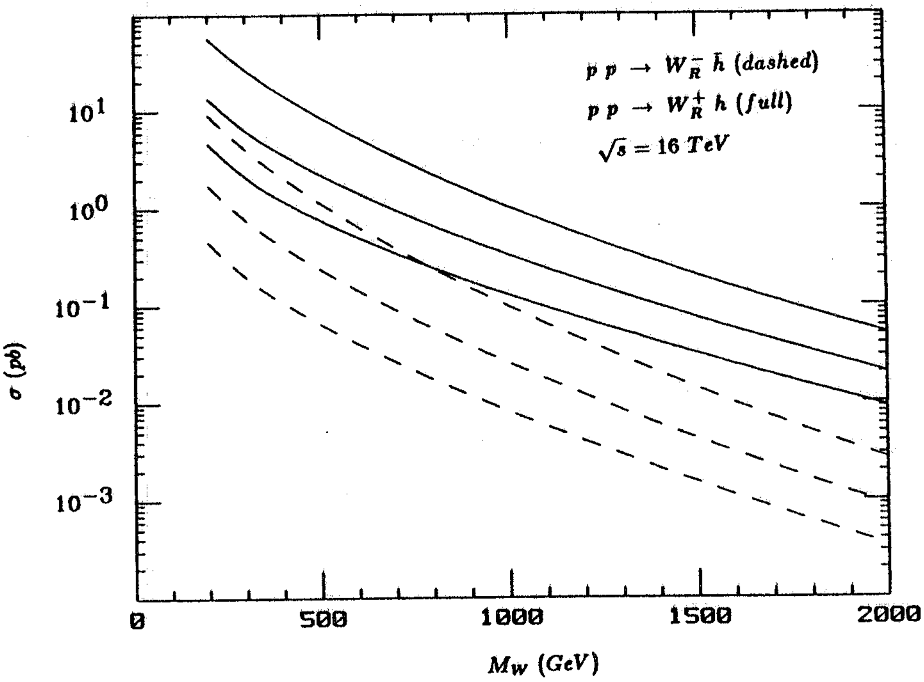


Fig. 1: Cross-sections at LHC for $pp \to W_R^+(full)\ h$ and $pp \to W_R^-(dashed)\ \bar{h}$ as a function of M_{W_R}. The three different curves, from the upper to the lower curve, correspond to $m_h = 0.3,\ 0.6,\ 0.9\ TeV$, respectively.

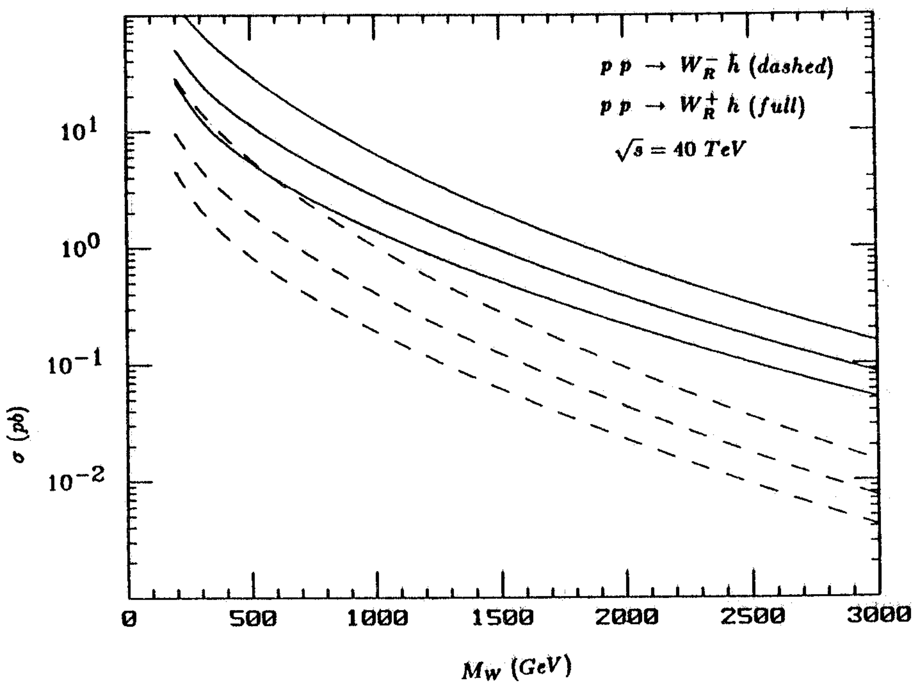


Fig. 2: Same as Fig. 1, except at SSC.

Table 2: Discovery limits for W_R and W_I at LHC and SSC.

	W_R^+	W_R^-	W_I	W_I^+
LHC	2 - 2.5 TeV	1 - 1.5 TeV	1.5 - 2.2 TeV	1 - 1.5 TeV
SSC	2.5 - 3 TeV	1.2 - 2 TeV	2 - 2.5 TeV	1.2 - 2 TeV

Next, we turn our attention to the possible final state signatures. The decay modes of the leptoquark h depend on the superpotential. If the $N_e{}^c$ in Table 1 is given negative R-parity, then the possible final state signatures are [13]

$$a)\ jet\ +\ l^+l^-\ +\ \not{p}_T\ ,$$

$$b)\ jet\ +\ e^-\ +\ \not{p}_T\ .$$

These are obtained from the decay modes

$$a)\ h \to d\ +\ \tilde{\nu}\ ,$$

$$b)\ h \to u\ +\ \tilde{e}^-\ ,$$

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 $h \to d\ +\ \tilde{N}_e$ 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_E by at least one order of magnitude [3,4]. Although there are no direct experimental constraints on the W_R mass, one can still put an indirect limit using the Z' mass bound. Hence, by purely kinematical reasons, owing to phase-space suppression, the largely dominant decay mode is expected to be the one involving n, instead of h, E, N_E

$$W_R{}^+ \to e_L{}^+\ +\ n_R\ .$$

The estimated branching ratio for this mode is larger than 10%. This yields the possible final state signatures

$$a')\ W_R{}^+ \rightarrow e^+ \ +\ \gamma \ +\ \not{p}_T(\tilde{\gamma})\ ,$$

if the LSP is, for example, the photino $\tilde{\gamma}$ and the n decays before leaving the detector, or

$$b')\ W_R{}^+ \rightarrow e^+ \ +\ \not{p}_T(n)\ ,$$

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 $\tilde{\gamma}$ and the remaining neutralinos is needed.

The partial widths for the W_R-decays can be expressed in terms of the ratio between the W_R and W_L mass. Denoting by n_g the number of generations of heavy exotic fermions h, E, N_E, one has [14]

$$\Gamma(\ W_R \rightarrow fermions\ ) \ = \ (\ 0.69\ +\ 1.15\ n_g\ )\ GeV \times \left[\frac{M_{W_R}}{M_{W_L}} \right] . \qquad (10)$$

Neglecting phase space effects, one also finds [14]

$$\Gamma(\ W_R \rightarrow gauge\ bosons\ and\ Higgs\ ) \ = \ 0.23\ GeV \times \left[\frac{M_{W_R}}{M_{W_L}} \right] , \qquad (11)$$

as well as

$$\Gamma(\ W_R \rightarrow SUSY\ partners\ ) \ = \ 1.74\ GeV \times \left[\frac{M_{W_R}}{M_{W_L}} \right] . \qquad (12)$$

Taking $M_{W_R} = 2\ TeV$, together with the experimental value $M_{W_L} = 80\ GeV$, one gets $\Gamma_{tot}(W_R) = 70\ (150)\ GeV$ for $n_g = 0\ (3)$. Combining the h-decay and the W_R-decay gives rise to a final state with a very large invariant mass.

4. - Production of flavor-changing neutral gauge bosons

Recently, there has been some interest for the flavor-changing gauge boson W_I arising from the $SU(2)_L \times U(1)_Y \times SU(2)_I$ model [2]. This low-energy gauge group can arise when breaking E_6, and the generators of $SU(2)_I$ commute with the electric charge [1]. The pair of conjugate non-hermitian gauge bosons denoted by W_I and $W_I{}^\dagger$ correspond to the non-diagonal generators of $SU(2)_I$. W_I has negative R-parity and non-zero lepton number $L = -1$. It couples to fermions according to the lagrangian

$$\mathcal{L} = \frac{1}{\sqrt{2}} g W_I{}^{\mu} \Big(\bar{h} \gamma_\mu d_R + \bar{e} \gamma_\mu E_L + \bar{\nu} \gamma_\mu (N_E)_L + \bar{\nu^c} \gamma_\mu n^c{}_L \Big) + h.c. , \quad (13)$$

where g is the usual $SU(2)_L$ coupling constant and it has been assumed that both $SU(2)$ factors of the low-energy group originate at a common scale from the breakdown of a larger group.

The dominant W_I production mechanisms are $g + d \to h + W_I$ and $g + \bar{d} \to \bar{h} + W_I{}^{\dagger}$. These parton-level processes give rise to the spin and color averaged matrix element

$$< |\mathcal{M}_I|^2 > = \frac{G_F M_W{}^2}{3\sqrt{2}} 4\pi\alpha_s \Bigg[- \left(\frac{t'}{s} + \frac{s}{t'} \right) \left(2 + \frac{m_h{}^2}{M_{W_I}{}^2} \right) - 2\frac{m_h{}^2}{M_{W_I}{}^2}$$
$$+ 2 \left(2M_{W_I}{}^2 - m_h{}^2 - \frac{m_h{}^4}{M_{W_I}{}^2} \right) \left(\frac{1}{s} + \frac{1}{t'} \right)$$
$$+ \frac{2}{st'} \left(- \frac{m_h{}^6}{M_{W_I}{}^2} + 3 m_h{}^2 M_{W_I}{}^2 - 2 M_{W_I}{}^4 \right) \Bigg] , \quad (14)$$

where $t' = t - m_h{}^2$. Comparison with eq. (3) shows that the percentage difference between the production cross-sections for W_I and W_R is generally small, especially if the mass of the leptoquark h and the mass of the W_R gauge boson have values of the same order of magnitude. 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 confirmed by the explicit numerical results in Figs. 3,4. Comparison with Figs. 1,2 also confirms the similarity of the results for W_I production with the cross-section for W_R.

Once again the production of flavor-changing neutral gauge bosons (Figs. 3,4) yields numerical estimates of the same order of magnitude at LHC and SSC, owing to the different luminosities for the two hadron colliders. Note that our numerical results do not agree with those reported in Figs. 2, 3, 5, 6 of ref. [2]. In the r.h.s. of eq. (5) of ref. [2] a factor 4 is missing, with respect to the production cross-section obtained from our eqs. (2,14). This is also repeated in eq. (2.95) of ref. [14]. The discrepancy approximately accounts for the difference in the numerical results obtained in refs. [2,14].

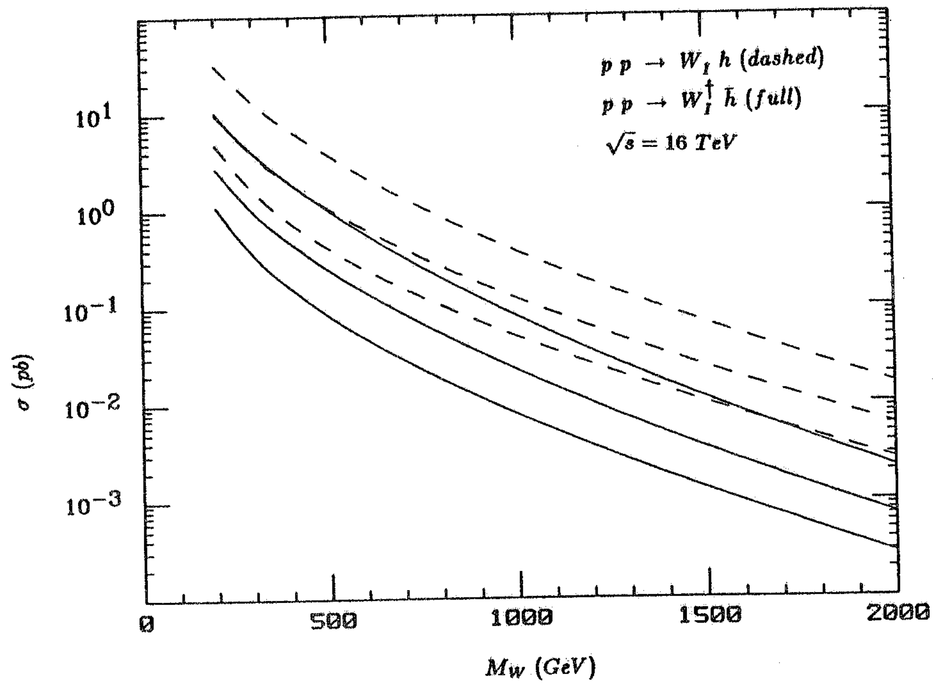


Fig. 3: Same as Fig. 1, except for W_I(*dashed*) h and $W_I^\dagger$(*full*) $\bar{h}$.

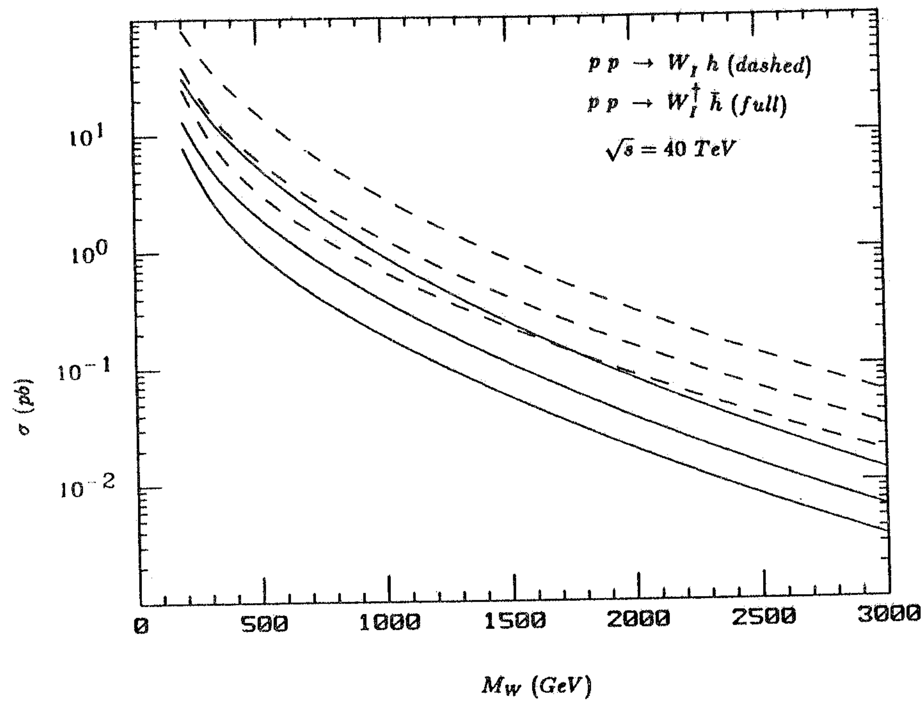


Fig. 4: Same as Fig. 3, except at SSC.

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 as in sec. 3, 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%, Figs. 3,4 yield discovery limits for the flavor-changing gauge bosons, also summarized in

Table 2. We find that, at both hadron colliders, W_I will be observable up to a mass of 1.5-2 TeV, whereas for $W_I{}^\dagger$ the discovery limit is given by a mass of about 1-1.5 TeV, depending on the leptoquark mass.

References

[1] D. London and J.L. Rosner, Phys. Rev. D34 (1986) 1530; J.L. Rosner, Comments Nucl. Part. Phys. 15 (1986) 195.

[2] T.G. Rizzo, Phys. Rev. D38 (1988) 71.

[3] E. Ma, Phys. Rev. D36 (1987) 274; D.S. Babu, X-G. He and E. Ma, ibid. 878.

[4] J.F. Gunion, J.L Hewett, E. Ma and T.G. Rizzo, Int. J. Mod. Phys. A2 (1987) 1199.

[5] R.N. Mohapatra, Unification and Supersymmetry (Springer, New York, 1986) and references therein; Q. Shafi, and Ch. Wetterich, Phys. Lett. 73B (1978) 65; V. Elias, J.C. Pati and A. Salam, Phys. Lett. 73B (1978) 451.

[6] G. Beall, M. Bander and A. Soni, Phys. Rev. Lett. 48 (1982) 848.

[7] G. Ecker and W. Grimus, Nucl. Phys. B258 (1985) 328.

[8] P. Colangelo and G. Nardulli, Universitá di Bari preprint BARI TH/90-67 (1990).

[9] J. Carr et al., Phys. Rev. Lett. 51 (1983) 627; D.P. Stoker et al., Phys. Rev. Lett. 54 (1985) 1887.

[10] P. Langacker and S. Uma Sankar, Phys. Rev. D40 (1989) 1569.

[11] M. Diemoz, F. Ferroni, E. Longo and G. Martinelli, Z. Phys. C39 (1988) 27.

[12] D. Duke and J. Owens, Phys. Rev. D30 (1984) 49.

[13] V. Barger, N.G. Deshpande and J.F. Gunion, in Proceedings of the Summer Study on the Physics of the Superconducting Super Collider, Snowmass, Colorado, 1986, eds. R. Donaldson and J. Marx (New York, 1987)

[14] J.L. Hewett and T.G. Rizzo, Phys. Rep. 183 (1989) 193.