CERN-TH.4876/87

HIGH-ENERGY TESTS FOR A POSSIBLE STRONG SECTOR IN THE STANDARD MODEL

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ABSTRACT

We discuss some high-energy tests of an electroweak model which can be thought of as an effective description of the standard model when the Higgs sector becomes strongly interacting. The model contains, in addition to the standard gauge bosons, three new vector particles arising as bound states due to the strong interaction in the Higgs sector. Precise measurements of m_W and m_Z , and asymmetries at SLC and LEP can give strong constraints on the couplings of these new vector bosons to the fermions and on the mixing angles to standard W and Z. We discuss how the new vector particles could provide for an explanation of the $\nu\nu jj$ events seen at the collider, and how they may be seen at the Tevatron, at LHC and at SSC.

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1. - INTRODUCTION

Exciting prospects exist at this time for extensive experimentation at new high-energy accelerators which are under construction or are being designed, and at present machines after detector and intensity improvements. Such prospects make it important at this time to develop schemes for possible deviations from the standard electroweak model which are physically reasonable and also sufficiently simple to be used as possible backgrounds for its quantitative tests. The model [1] we shall use in this paper incorporates what is expected to be a most visible feature of a strong interacting electroweak sector, namely the appearance of new vector resonances which would mix with W and Z and so give rise to visible distortions in the low-energy and medium-energy phenomenology. On the other hand, this model [1] has an independent logical basis, since it can be seen as an alternative way of realizing the electroweak symmetry breaking, namely in a non-linear way rather than in the usual linear form. One way of looking at the standard electroweak symmetry breaking is in fact to concentrate on its scalar sector, which has an initial $SU(2)_L \otimes SU(2)_R$ symmetry which breaks down spontaneously into a diagonal $SU(2)_D$. The Higgs field initially belongs to $(\frac{1}{2}, \frac{1}{2})$ of $SU(2)_L \otimes SU(2)_R$. After breaking, it behaves as $\underline{1} \oplus \underline{3}$ under the unbroken $SU(2)_D$: the singlet is the physical Higgs, the triplets are the Goldstones, which are later absorbed to give masses to W and Z. Now, such a diagonal $SU(2)_D$ which is left unbroken is indeed the so-called "custodial $SU(2)$ ", many times invoked in technicolour theories, which has a rôle of guaranteeing $\rho = 1$ apart from weak corrections. In fact, the linear realization of such a symmetry-breaking scheme is exactly the standard model in its usual form. From these considerations, one is naturally led to speculate about a possible non-linear realization (of the type of a non-linear σ -model). Such a non-linear realization, treated at the classical level, reproduces the limit of the large Higgs mass, $m_H \rightarrow \infty$, of the standard linear realization (m_H is the Higgs mass appearing in the Lagrangian). Effectively the non-linear condition on the field has resulted in eliminating one scalar degree of freedom, so that the model no longer possesses in its classical Lagrangian the degree of freedom corresponding to the physical Higgs particle. Such a feature is considered welcome, because of the difficulties encountered with a Higgs particle, and is one of the main motivations for our choice of the non-linear realization of symmetry breaking. It is easy to convince oneself that it is quite generally impossible, in the linear approach to symmetry breaking, to absorb all scalar degrees of freedom through the Higgs phenomenon so that no physical scalars remain. In the non-linear approach, this is indeed possible. One knows that for $m_H \rightarrow \infty$ one tends towards a strong-interaction limit [2,3] so that such a situation plausibly appears as one where a strong-interaction sector develops inside the electroweak theory.

The analogy with QCD is evident. In QCD with N massless quarks, the symmetry breaking is from $SU(N)_L \otimes SU(N)_R \otimes U(1)$ to $U(N)_D$, while the Goldstones are the

known pseudoscalars which remain present in the physical spectrum, and, furthermore, massive vector bound states are formed. In our case, the Goldstones are absorbed in giving masses to W and Z, but we must still expect massive vector bound states to appear. Since they will mix with W and Z, they are of primary phenomenological importance. A convenient way to construct an effective Lagrangian also including the vector bound states is simply to extend the initial $SU(2)_L \otimes SU(2)_R$ group into an $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V$ and construct a non-linear breaking into a diagonal $SU(2)_D$. In such a non-linear scheme, the initial Goldstones will correspond to a total of six degrees of freedom, corresponding to the six generators in the quotient $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V / SU(2)_D$. Three of them will give masses to W and Z after $SU(2)_L \otimes U(1)$ is gauged; the other three will give masses to an $SU(2)$ triplet of gauge bosons, to be called V^+ , V^0 , V^- . The counting shows that all the scalar degrees of freedom have disappeared, to give masses through the Higgs mechanism, and there is no room for surviving physical Higgs particles among the Lagrangian degrees of freedom. Such a counting, leaving zero surviving scalars, would not have been possible in a linear realization of symmetry breaking.

This is the main outline of the model. Because of its maximal simplicity, it can be seen as nothing other than a convenient way of offering a background against which to test the validity of the standard electroweak scheme. However, if one has faith that something different will sooner or later have to appear, one may also take the model as a relevant physical alternative. Independently of the attitude one wants to take, in this paper we will only be concerned with the phenomenology. We shall take the quantitative predictions of the model and test them against existing or future experiments. The low-energy phenomenology has already been tested and has resulted in bounds for the parameters of the model [4,5] which will be taken into account here.

Concerning the high-energy phenomenology, we devote a part of our paper to the study of the constraints on the model considered here, coming from experiments which will be performed at the future e^+e^- colliders, LEP and SLC. We expect that a set of precision measurements could test the reliability of the standard model and put limitations on its possible extensions. We are particularly interested in this last possibility and we will adopt the following point of view. Given certain experimental accuracies for observables measured at LEP and SLC, and provided that no deviations from the standard model are observed, we will obtain limits on the parameters which characterize the model. The quantities we have selected in order to perform such an analysis are: the vector boson masses m_W and m_Z , the partial width of the Z into charged leptons, the forward-backward asymmetry A_{FB} (both without and with longitudinally-polarized initial electron beam) and the left-right asymmetry A_{LR} .

Concerning existing high-energy experiments, the anomalous events of UAl [6], although still needing experimental confirmation, are the first ones which come to mind. If established, such an anomaly would require some not obvious modification of the standard model. All attempts at explanation have so far failed. We shall see that the model considered here would naturally account for such events; however, we repeat, additional experimental evidence is needed. Independently of the possibility of explaining the two UAl anomalous events as the decay of the new vector states from the strong interacting Higgs, it is anyway interesting to investigate the production of such particles at future hadron colliders and compare it with the standard model WW and WZ backgrounds.

In Section 2, we describe the model. In Section 3, we recall the definitions and the standard model predictions for the relevant observables in the high-energy phenomenology, as well as their expressions derived in the present model. Computer programs were set up to obtain the above-mentioned limitations in the parameter space and the main results are presented in Section 4. In Section 5, we show how our model can explain the anomalous UAl events and present the expectations for future proton colliders.

2. - EFFECTIVE LAGRANGIAN DESCRIPTION. POSSIBLE VECTOR RESONANCES

We shall review here the effective Lagrangian description [1] of the longitudinal scalar sector of the standard model which should apply to the case when it becomes strong-interacting. Independently of its dynamical origin, one may also consider the resulting effective Lagrangian as a phenomenological way of parametrizing possible deviations from the standard model in order to assess quantitatively its limits of validity. In this sense, the determination from data of bounds on the parameters of the effective Lagrangian might, in a way, play a similar rôle to the cut-off limits in terms of which theories (for instance QED) have sometimes been tested. The test would, in this case, be carried out against a background deviation resulting from a strong-interacting electroweak sector, of which our effective Lagrangian contains some simple and general features independent of its dynamical origin.

The symmetry-breaking (SB) sector of the standard model can be studied through the help of an effective Lagrangian. In fact, the SB is completely determined once the original invariance group of the theory, G , and the residual invariance group, H (invariance of the vacuum), are given. The Goldstone bosons are described as co-ordinates of the quotient space G/H , and indeed their quantum numbers are the same as for the generators of G/H . What we need in our particular problem are at least three Goldstone bosons to give masses to $W^\pm$ and Z , and that $H \supset U(1)_{e.m.}$.

Further information comes from the experimental result $\rho \approx 1$, which can be shown to hold automatically (apart from weak corrections) if the theory contains an unbroken $SU(2)$ ["custodial" $SU(2)$]. This implies that $H \supset SU(2)$. The minimal model emerging from these considerations is characterized by $G = SU(2) \otimes SU(2)$ and $H = SU(2)$. The standard model is just a linear realization of this case, with the Higgs field belonging to the representation $(\frac{1}{2}, \frac{1}{2})$ of $SU(2) \otimes SU(2)$ and a Higgs potential which is $SU(2) \otimes SU(2)$ symmetric. The SB leaves the diagonal $SU(2)$ as a symmetry of the vacuum, and the Higgs fields behaves as $\underline{1} \oplus \underline{3}$ under this $SU(2)$. The singlet remains as the physical Higgs, whereas the fields of the triplet are the Goldstone bosons which are eaten up by the vectors. For comparison, let us recall that in QCD with N massless quarks, $G = SU(N) \otimes SU(N) \otimes U(1)$ and $H = U(N)$.

Clearly, if we have to think of the Higgs description as an effective one, then it is more economical to use a non-linear realization which contains no physical Higgs but only Goldstone particles. This non-linear realization is equivalent to the Higgs description in the limit $m_H \rightarrow \infty$ (at least at one-loop level [2,3]). In fact, let us consider the Lagrangian of the Higgs sector:

$$\mathcal{L}_{\text{Higgs}} = \frac{1}{4} \text{Tr}(\partial_\mu M \partial^\mu M^\dagger) - \frac{m_H^2}{8f^2} \left(\frac{1}{2} \text{Tr}(MM^\dagger) - f^2 \right)^2 \quad (2.1)$$

where we have described the Higgs field in terms of a 2×2 matrix M :

$$M = \sqrt{2} \begin{pmatrix} \varphi_0 & -\varphi_-^* \\ \varphi_- & \varphi_0^* \end{pmatrix} \quad (2.2)$$

By performing a polar decomposition of the matrix M

$$M = \rho U \quad (2.3)$$

with

$$\rho = \sqrt{2} (|\varphi_0|^2 + |\varphi_-|^2)^{1/2}, \quad U^\dagger U = 1 \quad (2.4)$$

we get

$$\mathcal{L}_{\text{Higgs}} = \frac{1}{4} \partial_\mu \rho \partial^\mu \rho + \frac{1}{4} \rho^2 \text{Tr}(\partial_\mu U \partial^\mu U^\dagger) - \frac{m_H^2}{8f^2} (\rho^2 - f^2)^2 \quad (2.5)$$

In the limit $m_H \rightarrow \infty$, we have $\rho \rightarrow f$ and

$$\mathcal{L}_{\text{Higgs}} \rightarrow \mathcal{L} = \frac{f^2}{4} \text{Tr}(\partial_\mu U \partial^\mu U^\dagger) \quad (2.6)$$

$\mathcal{L}$ is the Lagrangian of a non-linear σ -model described by the unitary matrix U . Of course, the physical degrees of freedom are only three, due to the non-linear constraint (2.4). Notice that $\mathcal{L}$ is invariant under the transformations of $SU(2)_L \otimes SU(2)_R$ given by

$$U \rightarrow g_L U g_R^\dagger \quad g_L \in SU(2)_L, \quad g_R \in SU(2)_R \quad (2.7)$$

The non-linear condition (2.4) realizes the breaking down to the diagonal subgroup $SU(2)$, because one must have

$$\langle U \rangle_0 \propto 1 \quad (2.8)$$

We are guaranteed that in the low-energy regime and in the limit considered, the Lagrangian (2.6) gives rise to the same amplitudes that we get in the linear model. Analogous considerations are used in QCD, where one describes the sector of the pseudoscalar mesons (the Goldstone particles of the problem) by a non-linear σ -model analogous to the previous one (changing the group, of course). As is well known, the use of this σ -model in hadronic physics is very useful, because one can derive a number of phenomenological relations without knowing the underlying strong interacting theory (QCD in this particular instance). Also, in hadronic physics one knows that vector resonances are formed. The typical vector resonance is the ρ particle which has been conjectured to behave as a massive Yang-Mills particle. We can try to apply this idea to our case. We can describe the possible vector resonances as Yang-Mills particles, receiving mass from the same SB which gives mass to the standard gauge bosons. Of course, we will need more Goldstone bosons for this game. One of the simplest possibilities (corresponding to the so-called "hidden gauge symmetry" [7]) is to have an original group G given by $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V$ broken down again to $SU(2)_D$. The first two groups in G are the same as before, whereas the extra $SU(2)_V$ will allow us to describe three new massive vector bosons. In fact, the quotient space G/H will give us six Goldstone bosons, three of which will give mass to $W^\pm$ and Z , whereas the other three will give mass to the new vector particles $V^\pm$ and V^0 . In analogy to what we have done before, we will introduce here the Goldstone bosons through two 2×2 unitary matrices $L(x)$ and $R(x)$:

$$L^\dagger L = R^\dagger R = 1 \quad (2.9)$$

transforming as

$$L(x) \rightarrow g_L L(x) g, \quad R(x) \rightarrow g_R R(x) g$$

$$g_L \in SU(2)_L, \quad g_R \in SU(2)_R, \quad g \in SU(2)_V \quad (2.10)$$

Notice that if for a moment we forget about the non-linear constraints (2.9), then we can say that L and R belong to the following representations of G:

$$L(x) \in \left(\frac{1}{2}, 0, \frac{1}{2}\right), \quad R(x) \in \left(0, \frac{1}{2}, \frac{1}{2}\right) \quad (2.11)$$

The breaking of $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V$ down to $SU(2)_D$ is again realized through the requirement that the expectation value of the fields preserves the non-linear constraints (2.9) [see also Eq. (2.8)]:

$$\langle L(x) \rangle_0 \propto 1, \quad \langle R(x) \rangle_0 \propto 1 \quad (2.12)$$

We are now in the position of writing down a Lagrangian invariant under G, as a function of the non-linear fields L(x) and R(x). We observe that due to the non-linear constraints and the transformation properties of the fields, no non-trivial terms without derivatives exist. If furthermore we require a discrete symmetry under the exchange $L(x) \leftrightarrow R(x)$ (which will be justified below), then the most general Lagrangian containing at most two derivatives is:

$$\mathcal{L} = \frac{1+\alpha}{4} f^2 [\text{Tr}(\partial_\mu L^\dagger \partial^\mu L) + \text{Tr}(\partial_\mu R^\dagger \partial^\mu R)] + \frac{1-\alpha}{2} f^2 \text{Tr}(L^\dagger \partial_\mu L R^\dagger \partial^\mu R) \quad (2.13)$$

Here α is an arbitrary parameter, whereas f is the vev of the Higgs field. In fact, if we take $\alpha = 0$ in this Lagrangian, we recover the expression (2.6) for the non-linear realization of the standard model (in the following, a more interesting limit will appear in which one gets the same result). This can be seen by using the identification $U = LR^\dagger$ and the relation:

$$\text{Tr}(\partial_\mu U^\dagger \partial^\mu U) = \text{Tr}(\partial_\mu L^\dagger \partial^\mu L + \partial_\mu R^\dagger \partial^\mu R + 2L^\dagger \partial_\mu L R^\dagger \partial^\mu R) \quad (2.14)$$

Notice that without the requirement of symmetry under $L \leftrightarrow R$, we could not have obtained the non-linear realization of the standard model in this limit. In fact, in this last case we have invariance under $U \leftrightarrow U^\dagger$, which corresponds to $L \leftrightarrow R$.

$SU(2)_V$ is called a hidden gauge symmetry because U is a singlet under it, and

therefore this extra symmetry cannot be observed as long as we deal with U only. The Lagrangian (2.13) can be rewritten in the form:

$$\mathcal{L} = -\frac{1}{4}f^2 \left[\text{Tr} (L^\dagger \partial_\mu L - R^\dagger \partial_\mu R)^2 + 2 \text{Tr} (L^\dagger \partial_\mu L + R^\dagger \partial_\mu R)^2 \right] \quad (2.15)$$

The gauging of this Lagrangian can now be made in a trivial way by modifying the derivatives into the covariant derivatives with respect to the gauge group $SU(2)_L \otimes U(1)_R \otimes SU(2)_V$:

$$\begin{aligned} \partial_\mu L &\rightarrow D_\mu^{(L)} L = (\partial_\mu + W_\mu^{(0)}) L - L V_\mu^{(0)} \equiv \nabla_\mu^{(L)} L - L V_\mu^{(0)} \\ \partial_\mu R &\rightarrow D_\mu^{(R)} R = (\partial_\mu + Y_\mu) R - R V_\mu^{(0)} \equiv \nabla_\mu^{(R)} R - R V_\mu^{(0)} \end{aligned} \quad (2.16)$$

Here we have used a compact matrix notation for the vector fields:

$$Y_\mu = i Y_\mu \frac{\tau_3}{2}, \quad W_\mu^{(0)} = i \vec{W}_\mu^{(0)} \cdot \frac{\vec{\tau}}{2}, \quad V_\mu^{(0)} = i \vec{V}_\mu^{(0)} \cdot \frac{\vec{\tau}}{2} \quad (2.17)$$

and the superscript zero for W and V has been added to allow for a simpler notation after the fields get mixed. Then we get

$$\mathcal{L} = -\frac{1}{4}f^2 \left[\text{Tr} (L^\dagger \nabla_\mu^{(L)} L - R^\dagger \nabla_\mu^{(R)} R)^2 + 2 \text{Tr} (L^\dagger \nabla_\mu^{(L)} L + R^\dagger \nabla_\mu^{(R)} R - 2V_\mu^{(0)})^2 \right] \quad (2.18)$$

The coupling constants have been inserted into the definition of the vector fields. To this Lagrangian we must add the kinetic terms for the vector fields

$$\begin{aligned} \mathcal{L}_{\text{kin. terms}} &= -\frac{1}{4g^2} \text{Tr} [F_{\mu\nu}(W^{(0)}) F^{\mu\nu}(W^{(0)})] - \frac{1}{4g'^2} \text{Tr} [F_{\mu\nu}(Y) F^{\mu\nu}(Y)] \\ &\quad - \frac{1}{g''^2} \text{Tr} [F_{\mu\nu}(V^{(0)}) F^{\mu\nu}(V^{(0)})] \end{aligned} \quad (2.19)$$

From these expressions we see that in the strong coupling limit, $g'' \rightarrow \infty$, the kinetic term for the field V is negligible and the field V itself becomes an auxiliary field which can be eliminated through its equation of motion coming from (2.18):

$$V_\mu^{(0)} = \frac{1}{2} (L^\dagger \nabla_\mu^{(L)} L + R^\dagger \nabla_\mu^{(R)} R) \quad (2.20)$$

Using this relation, (2.18) reduces to the first term, which does not depend on α , and we recover the non-linear formulation of the standard model, coupled to the Weinberg-Salam gauge bosons.

Going into a unitary gauge ($L=1$, $R=1$), we can read directly from (2.18) the mass matrix of the gauge vector bosons. The unitary gauge can be reached by means of the gauge transformation:

$$W_\mu^{(0)} = U^\dagger \tilde{W} U + U^\dagger \partial_\mu U, \quad V_\mu^{(0)} = R^\dagger \tilde{V} R + R^\dagger \partial_\mu R \quad (2.21)$$

where $U = LR^\dagger$. Then we get

$$\mathcal{L}_{\text{tot}} = -\frac{1}{4} [\text{Tr}(\tilde{W} - Y)^2 + \alpha \text{Tr}(\tilde{W} + Y - 2\tilde{V})^2] \quad (2.22)$$

+ kinetic terms for the gauge fields

Finally we perform a rescaling of the fields according $\tilde{W} \rightarrow g\tilde{W}$, $Y \rightarrow g'Y$, $2\tilde{V} \rightarrow g''\tilde{V}$, and after a diagonalization of the charged and the neutral sectors, we get the physical vector bosons $W^\pm$, $V^\pm$, A , Z^0 and V^0 .

For the charged sector we have

$$\begin{pmatrix} W_i \\ V_i \end{pmatrix} = \begin{pmatrix} \cos \varphi & -\sin \varphi \\ \sin \varphi & \cos \varphi \end{pmatrix} \begin{pmatrix} \tilde{W}_i \\ \tilde{V}_i \end{pmatrix} \quad (i=1,2) \quad (2.23)$$

The mass matrix diagonalization can be performed exactly. However, we are mostly interested in the limit $g'' \gg g$. In this limit, the mixing angle and the masses are given by:

$$\varphi = -\frac{g}{g''} \frac{1}{1 - \left(\frac{\tilde{m}_W}{m_V}\right)^2}$$

$$m_W^2 = \frac{1}{4} f^2 g^2 \left(1 - \frac{g^2}{g'^2} \frac{1}{1 - \left(\frac{\bar{m}_W}{m_V} \right)^2} \right)$$

$$m_V^2 = \frac{1}{4} f^2 g'^2 \quad (2.24)$$

For the neutral sector we have

$$\begin{pmatrix} A \\ Z^0 \\ V^0 \end{pmatrix} = \begin{pmatrix} 0 & \cos \psi & \sin \psi \\ \cos \xi & \sin \xi \sin \psi & -\sin \xi \cos \psi \\ \sin \xi & -\cos \xi \sin \psi & \cos \xi \cos \psi \end{pmatrix} \begin{pmatrix} g/G & -g'/G & 0 \\ g'/G & g/G & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} \tilde{W}_3 \\ Y \\ \tilde{V}_3 \end{pmatrix} \quad (2.25)$$

where $G = \sqrt{g^2 + g'^2}$. The mixing angles and masses are given (in the same limit) by

$$\xi = - \frac{g^2 + g'^2}{g' G} \cdot \frac{1}{1 - \left(\frac{\bar{m}_W}{m_V} \right)^2 \frac{1}{\cos^2 \theta}}, \quad \psi = \frac{2 g g'}{g'' G}$$

$$m_A^2 = 0, \quad m_Z^2 = \frac{1}{4} f^2 (g^2 + g'^2) \left[1 - \frac{(g^2 - g'^2)^2}{G^2 g'^2} \frac{1}{1 - \left(\frac{\bar{m}_W}{m_V} \right)^2 \frac{1}{\cos^2 \theta}} \right]$$

$$m_{V^0}^2 = \frac{1}{4} f^2 g'^2 \quad (2.26)$$

where

$$\sin \theta = \frac{g'}{\sqrt{g^2 + g'^2}} \quad (2.27)$$

In these formulae we have used the exact expressions of Ref. [1] taken at the order $(g/g'')^2$ with $(\bar{m}_W/m_V)^2$ fixed, where $\bar{m}_W$ can be taken as the limiting value of m_W for $g'' \rightarrow \infty$.

For $g'' \rightarrow \infty$, we recover the standard model mass formulae. This is what we anticipated because, again, in this limit the kinetic term for the field V is negligible and we recover the gauged non-linear σ -model.

Let us now consider the fermions of the GWS model and denote them in a compact way by ψ_L and ψ_R , where

$$\psi_L \in \left(\frac{1}{2}, 0, 0\right) \quad , \quad \psi_R \in \left(0, \frac{1}{2}, 0\right) \quad (2.28)$$

under the group $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V$ (this cannot, of course, hold for the right-handed leptons if there are no right-handed neutrinos). Then the fermions couple to V only via the mixing with the W :

$$\begin{aligned} \mathcal{L}_{\text{fermion}} = & \bar{\psi}_L i \gamma^\mu \left(\partial_\mu + i \vec{W}_\mu^{(10)} \cdot \frac{\vec{\tau}}{2} + \frac{1}{2} i (B-L) Y_\mu \right) \psi_L \\ & + \bar{\psi}_R i \gamma^\mu \left(\partial_\mu + i \left(\frac{\tau_3}{2} + \frac{1}{2} (B-L) \right) Y_\mu \right) \psi_R \end{aligned} \quad (2.29)$$

from which we get the couplings

$$\begin{aligned} & (h_W W_\mu^i + h_V V_\mu^i) T_L^i \quad (i=1,2) \\ & Z_\mu^0 (A T_L^3 + B Q_{em}) + V_\mu^0 (C T_L^3 + D Q_{em}) + e Q_{em} A_\mu \end{aligned} \quad (2.30)$$

where

$$\begin{aligned} h_W &= g \cos \psi \quad , \quad h_V = g \sin \psi \quad , \quad e = \frac{g g'}{G} \cos \psi \\ A &= G \cos \zeta \quad , \quad B = -\frac{g'^2}{G} \cos \zeta \left(1 - \frac{g}{g'} \tan \zeta \sin \psi \right) \\ C &= G \sin \zeta \quad , \quad D = -\frac{g'^2}{G} \sin \zeta \left(1 + \frac{g}{g'} \cot \zeta \sin \psi \right) \end{aligned} \quad (2.31)$$

The charged couplings determine an effective current-current Hamiltonian whose strength by definition is $G_F \sqrt{2}$, for which one gets:

$$\frac{G_F}{\sqrt{2}} = \frac{1}{2f^2} \quad (2.32)$$

From the neutral coupling one gets the following effective neutral Hamiltonian:

$$\mathcal{H}_I^{m.c.} = \sqrt{2} G_F e [(T_3^L + X Q_{em})^2 + c Q_{em}^2] \quad (2.33)$$

where T_3^L and Q_{em} are respectively the third component of the isospin and the electromagnetic currents. Recall that in the standard model $\rho \sim 1$, X is $-\sin^2\theta_{WS}$ and $c = 0$. In the present case, we get:

$$\begin{aligned} e &= 1 \\ X &= -\sin^2\theta \left[1 + 2 \left(\frac{g}{g''} \right)^2 \cos 2\theta \right] \\ c &= 4 \left(\frac{g}{g''} \right)^2 \left(\frac{\bar{m}_W}{m_V} \right)^2 \sin^4\theta \end{aligned} \quad (2.34)$$

Also, from the electromagnetic coupling we get the expression for the electric charge:

$$e = g \sin\theta \left[1 - 2 \left(\frac{g}{g''} \right)^2 \sin^2\theta \right] \quad (2.35)$$

In order to fit the parameters (g, g', f) in the present model, one makes use of the previous expressions and of the experimental values for (G_F, e, X) . In particular one takes X as measured from the ratio

$$\frac{\sigma_{NC}(\nu_\mu N \rightarrow \nu_\mu X)}{\sigma_{CC}(\nu_\mu N \rightarrow \mu X)}$$

For this one can use a recent two-parameter fit obtained by Amaldi et al. and by Costa et al. [8]. This fit gives

$$\begin{aligned} \sin^2\theta &= 0.232 \pm 0.014 \pm [0.008] \\ e &= 0.999 \pm 0.013 \pm [0.008] \end{aligned} \quad (2.36)$$

where the errors inside the square brackets are due to the theoretical uncertainties. The parameters g'' , m_V (or g'', α) remain to be determined. Before doing so,

one has to discuss the possibility that the fermions have a direct coupling to the fields V . Although such couplings are absent in the original Lagrangian of the standard model, they can be present in the effective Lagrangian if they are allowed by the symmetry of the problem. In fact, such couplings can be derived by considering in correspondence of each doublet ϕ_L and ϕ_R new fermion fields

$$\chi_L = L^\dagger \psi_L, \quad \chi_R = R^\dagger \psi_R \quad (2.37)$$

and writing a general kinetic term

$$\begin{aligned} \mathcal{L}'_{\text{fermion}} = \mathcal{L}_{\text{fermion}} + b \bar{\chi}_L i \gamma^\mu (\partial_\mu + i \vec{V}_\mu^{(0)} \cdot \vec{\tau}) \frac{\bar{Z}}{2} + \frac{1}{2} i (B-L) \gamma_\mu \chi_L \\ + b' \bar{\chi}_R i \gamma^\mu (\partial_\mu + i \vec{V}_\mu^{(0)} \cdot \vec{\tau}) \frac{\bar{Z}}{2} + \frac{1}{2} i (B-L) \gamma_\mu \chi_R \end{aligned} \quad (2.38)$$

where $\mathcal{L}_{\text{fermion}}$ is given in Ep. (2.29) and b and b' are the unknown parameters of the effective Lagrangian. We would expect b to be of the order $g^2(g/g'')^2$ as arising from V-W mixing and one-loop corrections to the W-fermion vertex due to virtual emission of W's and Goldstone particles. From similar arguments b' is expected to be of the same order of b times an extra factor m_f/f (m_f is the fermion mass). The b' terms are absent in the leptonic sector if there are no right-handed neutrinos. In the hadronic sector they give rise to right-handed currents which contribute to K_L-K_S mass difference, and therefore from this difference we could derive an upper bound on b' . The bound on b' from this process and from the neutron dipole moment has been recently calculated by Neufeld [9]. He has found that these non-leptonic non-standard terms can explain CP violating observables like ϵ and ϵ' of the K-system when $2.6 \cdot 10^{-3} \leq |b'| \leq 8 \cdot 10^{-3}$, the upper bound coming from non-leptonic $\Delta S = 1$ weak decays. If b' is smaller than the previous limits, then CP-violating observables can be explained by allowing also for the standard contribution. The neutron electric dipole can also give results very near its present experimental bound.

For the sake of simplicity, we shall neglect the b' term in the following phenomenological analysis. It turns out useful to list also the vector and axial-vector couplings v_i and a_i ($i=1,2$) of the charged fermions to Z^0 ($i=1$) and V^0 ($i=2$):

$$v_1^f = -A T_3^L - 2B Q_{em}$$

$$\begin{aligned}
 a_1^f &= AT_3^L \\
 v_2^f &= -CT_3^L - 2DQ_{em} \\
 a_2^f &= CT_3^L
 \end{aligned}
 \tag{2.39}$$

$$A = G \cos \zeta (1+b)^{-1} \left[1 + b \sin^2 \theta \left(1 - \frac{\tan \zeta}{\tan \theta \sin \psi} \right) \right] \tag{2.40}$$

$$B = -G \sin^2 \theta \cos \zeta \left(1 - \frac{\tan \zeta \sin \psi}{\tan \theta} \right) \tag{2.41}$$

$$C = G \sin \zeta (1+b)^{-1} \left[1 + b \sin^2 \theta \left(1 + \frac{\cot \zeta}{\tan \theta \sin \psi} \right) \right] \tag{2.42}$$

$$D = -G \sin^2 \theta \sin \zeta \left(1 + \frac{\cot \zeta \sin \psi}{\tan \theta} \right) \tag{2.43}$$

$$\sqrt{g^2 + g'^2} = (4\sqrt{2} G_F m_Z^2)^{\frac{1}{2}} \left[1 + \frac{1}{2} \left(\frac{g}{g''} \right)^2 \frac{\cos^2 2\theta}{\cos^2 \theta - \frac{m_W^2}{m_V^2}} + \frac{b}{2} - \frac{b^2}{8} \frac{m_W^2}{m_V^2} \left(\frac{g''}{g} \right)^2 \right] \tag{2.44}$$

The quantities A,B,C,D are derived from Ref. [1]. $\sqrt{g^2 + g'^2}$ holds in the limit of small (g/g'') and b.

In the following, we will also need the trilinear gauge boson vertices. The couplings $V^0 W^+ W^-$ and $V^- W^+ Z^0$ coming from the kinetic terms are of the order g/g'' . However, higher derivative terms exist:

$$L_I = \frac{1}{2} g^{\mu\nu} F^{\mu\nu}(V^{(0)}) [L^+ D_\mu^{(L)} L - R^+ D_\mu^{(R)} R, L^+ D_\nu^{(L)} L - R^+ D_\nu^{(R)} R] \tag{2.45}$$

$$L_{II} = \frac{1}{2} g^{\mu\nu} F^{\mu\nu}(V^{(0)}) [L^+ D_\mu^{(L)} L + R^+ D_\mu^{(R)} R, L^+ D_\nu^{(L)} L + R^+ D_\nu^{(R)} R] \tag{2.46}$$

One can see that all the couplings coming from L_{II} are at least of the order $(g/g'')^2$ [1]. After the transformation (2.21) and the rescaling of the fields, we can rewrite L_I in the following way:

$$L_I = \frac{\gamma}{4} g'' \text{tr} F^{\mu\nu}(\tilde{V}) [g \tilde{W}_\mu - g' Y_\mu, g \tilde{W}_\nu - g' Y_\nu] \quad (2.47)$$

From this, we get the couplings

$$g_{V^0 W^+ W^-} = \frac{\gamma}{8} g'' g^2, \quad g_{V^\pm W^\mp Z} = \frac{\gamma}{8} g'' g G \quad (2.48)$$

where we have neglected $O(g/g'')$ terms.

3. - HIGH-ENERGY PREDICTION TESTS WITH e^+e^-

The next generation of e^+e^- colliders, reaching a centre-of-mass energy equal to m_Z , offers a unique opportunity to perform precision tests on the electroweak standard model (SM) and to constrain or discover any of its theoretical modifications.

Our strategy [10] will be:

- i) to select a number of quantities which, in the SM, are functions of m_Z and which are planned to be measured with high precision;
- ii) to compare the prediction of the SM and of the present model;
- iii) to constrain the parameter space of the model, by assuming that a certain experimental precision can be reached in measuring a given quantity, and that no discrepancies with the SM can be found.

We shall look for functions of m_Z essentially because m_Z is the quantity which will be measured with the largest precision at LEP I and SLC, and it will set the scale for all other high-precision tests. For instance, at LEP I [11]:

$$\delta m_Z = \pm 50 \text{ MeV} \quad (3.1)$$

(with a possible improvement of a factor 2-3 if a suitable absolute energy calibration, which is the main source of uncertainty, is achieved).

We shall now list the observables on which our analysis will be based:

i) $\Delta \equiv (m_W/m_Z)$

It can be directly measured at the improved CERN $p\bar{p}$ collider (ACOL). There one aims at a precision of [11]:

$$\delta\left(\frac{m_W}{m_Z}\right) = \pm 2 \cdot 10^{-3} \quad (3.2)$$

(This precision can be further improved by combining the results of LEP I/SLC and ACOL.)

In the SM, Δ coincides with $\cos\theta_W$ and can be thought of as a function of m_Z :

$$\Delta = \left\{ \frac{1}{2} + \left(\frac{1}{4} - \frac{\mu^2}{m_Z^2} \right)^{\frac{1}{2}} \right\}^{\frac{1}{2}} \quad (3.3)$$

$$\mu^2 = \frac{\mu_0^2}{(1 - \Delta r)} \quad (3.4)$$

$$\mu_0^2 = \frac{\pi \alpha(0)}{\sqrt{2} G_F} = (37.281 \text{ GeV})^2 \quad (3.5)$$

$$\frac{\alpha(0)}{(1 - \Delta r)} \sim \alpha(m_Z) \quad (3.6)$$

where Δr comes from standard model radiative corrections. We have taken $\Delta r = 0.0713$ assuming $m_t = 45 \text{ GeV}$ and $m_H = 100 \text{ GeV}$ [12].

ii) $\Gamma(Z \rightarrow \ell^+ \ell^-) \quad \ell = e, \mu$

The expected experimental precision we take as [11]

$$\frac{\delta \Gamma(Z \rightarrow \ell^+ \ell^-)}{\Gamma(Z \rightarrow \ell^+ \ell^-)} \sim 0.02 \quad (3.7)$$

(Also the total Z-width will be measured with the same precision [11] but we prefer $\Gamma(Z \rightarrow \ell^+ \ell^-)$, since it does not depend on the number of massless neutrino species or other possible, yet unknown, channels.)

In the standard model:

$$\Gamma(Z \rightarrow \ell^+ \ell^-) = \frac{G_F}{12\pi\sqrt{2}} m_Z^3 (1 - 4\hat{s}^2 + 8\hat{s}^4) (1 + \delta_e) \quad (3.8)$$

$$\hat{s}^2 = \frac{1}{2} - \left(\frac{1}{4} - \frac{m_Z^2}{m_Z^2} \right)^{\frac{1}{2}} \quad (3.9)$$

δ_e represents the electroweak radiative correction and turns out to be very small if the radiatively corrected formula (3.9) is used for $\hat{s}^2$ [13]:

$$\delta_e \leq 0.1\% \quad (3.10)$$

iii) $\underline{A_{FB}}$

$$A_{FB} = \frac{F - B}{F + B} \quad (3.11)$$

$$F \pm B = \left(\int_0^x \pm \int_{-x}^0 \right) dz \frac{d\sigma(e^+e^- \rightarrow f\bar{f})}{dz} \quad (3.12)$$

x is the detector acceptance ($x \leq 1$); z is the cosine of the angle between the fermion f and the beam direction. The most precise determination of A_{FB} can be performed on the Z peak and if the final state fermions are muons. For 40 days of running time with luminosity $\mathcal{L} = 10^{31} \text{ cm}^2\text{sec}^{-1}$, the experimental error is estimated to be [11]:

$$\delta A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}(m_Z) \sim 5 \cdot 10^{-3} \quad (3.13)$$

(At the LEP II energy of $\sqrt{s} = 190 \text{ GeV}$ and with $\int \mathcal{L} dt = 500 \text{ pb}^{-1}$, the experimental error is expected to be $\delta A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}(190 \text{ GeV}) \sim 0.03$ [14].)

In the SM, at $\sqrt{s} = m_Z$ and in the narrow-width approximation, one has [10]:

$$A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}(m_Z) = 2 \frac{x}{1 + \frac{x^2}{3}} \frac{v_e}{1 + v_e^2} \frac{2v_e + P_e(1 + v_e^2)}{1 + v_e^2 + 2P_e v_e} \quad (3.14)$$

where P_e is the degree of longitudinal polarization of the electron beam and $v_e = 1 - 4 \sin^2 \theta_W$ is proportional to the Z vector coupling constant for charged leptons. Observe that all terms containing P_e are proportional to v_e , which is small since $\sin^2 \theta_W$ is close to $1/4$. (This is true at all energies, not only at $\sqrt{s} = m_Z$.) Therefore the longitudinal polarization of e^- will affect the shape of

A_{FB} only slightly.

As far as the radiative corrections are concerned, A_{FB} in the SM can be written as the sum of three terms:

$$A_{FB} = A_{FB}^{\text{Born}} + \delta A_{FB}^{\text{QED}} + \delta A_{FB}^{\text{weak}} \quad (3.15)$$

The terms in (3.15) are not completely unambiguous. The QED radiative corrections consist of two parts: one is due to the real and virtual photon emission diagrams, which we identify with $\delta A_{FB}^{\text{QED}}$, following Ref. [11]. The other is due to the one-loop photon vacuum polarization diagrams and contributes mainly to the running of α . We take into account this contribution directly in A_{FB}^{Born} by running α or, equivalently, by using the radiatively-corrected formula (3.9) for $\sin^2\theta_W$. Finally, $\delta A_{FB}^{\text{weak}}$ are the purely radiative corrections to A_{FB} . At the resonance m_Z , $\delta A_{FB}^{\text{weak}}$ can vary by as much as ± 0.01 for $30 \text{ GeV} < m_t < 180 \text{ GeV}$, $10 \text{ GeV} < m_H < 1 \text{ TeV}$ and $90 \text{ GeV} < m_Z < 98 \text{ GeV}$ [15].

iv) A_{LR}

$$A_{LR} = \frac{\sigma_L - \sigma_R}{\sigma_L + \sigma_R} \quad (3.16)$$

where $\sigma_{L,R}$ are the cross-sections for $e_{L,R}^- + e^+ \rightarrow X$ and X can be any channel. At $\sqrt{s} = m_Z$, $\int \mathcal{L} dt = 40 \text{ pb}^{-1}$ and $P_e = 0.5$ (A_{LR} requires a longitudinally-polarized electron beam) the expected error is [10]:

$$\delta A_{LR}(m_Z) \sim 1.5 \cdot 10^{-3} \quad (3.17)$$

In the SM, at $\sqrt{s} = m_Z$, for a massless final fermion, in the narrow-width approximation, one has

$$A_{LR}(m_Z) = 2 P_e \frac{v_e}{1 + v_e^2} \quad (3.18)$$

Observe that in the SM, $A_{LR}(m_Z)$ is independent of the couplings of the final state. Similar considerations as for the A_{FB} apply to the radiative corrections of A_{LR} . In particular on top of the Z resonance, $\delta A_{LR}^{\text{weak}}$ can vary as much as ± 0.02 for $30 \text{ GeV} < m_t < 180 \text{ GeV}$, $10 \text{ GeV} < m_H < 1 \text{ TeV}$ and $90 \text{ GeV} < m_Z < 98 \text{ GeV}$ [15]. As far as the expressions of the various observables in the present model are concerned, they are easily derivable from Ref. [1].

For $\Delta = (m_W/m_Z)$ we obtain, always for small g/g'' ,

$$\Delta = \left[\frac{1}{2} + \sqrt{\frac{1}{4} - \frac{\mu^2}{m_Z^2} (1 + \varepsilon_f + \varepsilon_g + \varepsilon_\theta + \varepsilon_\Delta)^2} \right]^{\frac{1}{2}} \quad (3.19)$$

$$\varepsilon_f = -\frac{b}{2} + \frac{b^2}{8} \left(\frac{\bar{m}_W}{m_V} \right)^2 \left(\frac{g''}{g} \right)^2 \quad (3.20)$$

$$\varepsilon_g = \frac{b}{4} + \left(\frac{g}{g''} \right)^2 + \frac{b}{2} \left(\frac{\bar{m}_W}{m_V} \right)^2 \left(1 - \frac{b}{4} \left(\frac{g''}{g} \right)^2 \right) \quad (3.21)$$

$$\varepsilon_\theta = -\frac{b}{4} - \left(\frac{g}{g''} \right)^2 \cos 2\theta - \frac{b}{2} \left(\frac{\bar{m}_W}{m_V} \right)^2 \left(1 - \frac{b}{4} \left(\frac{g''}{g} \right)^2 \right) \quad (3.22)$$

$$\varepsilon_\Delta = \frac{1}{2} \left[\frac{\cos 2\theta}{\sin^2 \theta} \frac{1}{1 - \left(\frac{\bar{m}_W}{m_V} \right)^2} - \cot^2 \theta \frac{\cos^2 2\theta}{\cos^2 \theta - \left(\frac{\bar{m}_W}{m_V} \right)^2} \right] \left(\frac{g}{g''} \right)^2 \quad (3.23)$$

Formulae (3.19)-(3.23) must be compared to those of the standard model given in (3.3)-(3.6).

In order to discuss the other observables, we note that the parameter $\sin \theta = g'/\sqrt{g'^2 + g''^2}$ occurs repeatedly in our expressions and it can be conveniently expressed in the following form (for small b and g/g''):

$$\sin \theta = \sin \bar{\theta} (1 + \varepsilon_\theta) \quad (3.24)$$

$$\sin \bar{\theta} = \left[\frac{1}{2} - \sqrt{\frac{1}{4} - \frac{\mu^2}{m_Z^2} (1 + \varepsilon_Z)} \right]^{\frac{1}{2}} \quad (3.25)$$

$$\varepsilon_Z = 2\varepsilon_f + 2\varepsilon_g + 2\varepsilon_\theta \cot^2 \theta - \left(\frac{g}{g''} \right)^2 \frac{\cos^2 2\theta}{\cos^2 \theta} \frac{1}{\left[1 - \frac{1}{\cos^2 \theta} \left(\frac{\bar{m}_W}{m_V} \right)^2 \right]} \quad (3.26)$$

In formulae (3.20)-(3.26), for computational purposes, the experimental values for m_W and $\sin^2\theta_W$ can be used, since the ϵ_s are small quantities.

The partial width of Z in a charged, massless lepton pair is given by:

$$\Gamma(Z \rightarrow l^+ l^-) = \frac{m_Z}{48\pi} (v_l^2 + a_l^2) \quad (3.27)$$

to be compared with (3.8) and (3.9) of the SM. v_l and a_l were defined in Eq. (2.39).

$A_{FB}^{e^+e^- \rightarrow f\bar{f}}$ and $A_{LR}^{e^+e^- \rightarrow f\bar{f}}$ are defined in Eqs. (3.11), (3.12) and (3.16). Their expressions, for the general case where, in addition to the photon, two massive neutral vector bosons are exchanged, can be found, for instance, in Ref. [10]. Here we list them for convenience:

$$A_{FB}^{e^+e^- \rightarrow f\bar{f}} = \frac{x}{1 + \frac{x^2}{3}} \frac{(1-P_e) \sum_{h_f, h_e} h_f h_e |F(h_f, h_e)|^2 + 2P_e \sum_{h_f} h_f |F(h_f, 1)|^2}{(1-P_e) \sum_{h_f, h_e} |F(h_f, h_e)|^2 + 2P_e \sum_{h_f} |F(h_f, 1)|^2} \quad (3.28)$$

$$A_{LR}^{e^+e^- \rightarrow f\bar{f}} = P_e \frac{\sum_{h_f, h_e} h_e |F(h_f, h_e)|^2}{\sum_{h_f, h_e} |F(h_f, h_e)|^2} \quad (3.29)$$

where

$$F(h_f, h_e) = -\frac{e^2}{s} e_f + \frac{1}{4} \cdot \frac{(v_{1f} - h_f a_{1f})(v_{1e} - h_e a_{1e})}{s - m_1^2 + i m_1 \Gamma_1} + \frac{1}{4} \cdot \frac{(v_{2f} - h_f a_{2f})(v_{2e} - h_e a_{2e})}{s - m_2^2 + i m_2 \Gamma_2} \quad (3.30)$$

In the above formulae h_f, h_e ($= \pm 1$) are the helicities of f and e respectively; e_f is the electric charge of f (in units of e_{proton}); m_i and Γ_i are the mass and width of the neutral massive gauge bosons. The leptonic widths of Z and V^0 are given by:

$$\Gamma_1 = \left(24 + 40 \frac{B}{A} + 160 \frac{B^2}{A^2} \right) \Gamma_{1\nu\bar{\nu}} \quad (3.31)$$

$$\Gamma_{1\nu\bar{\nu}} = \frac{m_Z}{96\pi} A^2 \quad (3.32)$$

$$\Gamma_2 = 8 \left(3 + 6 \frac{D}{C} + 8 \frac{D^2}{C^2} \right) \Gamma_{2\nu\bar{\nu}} \quad (3.33)$$

$$\Gamma_{2\nu\bar{\nu}} = \frac{m_V}{96\pi} C^2 \quad (3.34)$$

In (3.33) we have not included a possible top contribution. Moreover, in the actual computations, we have assumed that Γ_2 is the total width of the V^0 , neglecting possible vector channels.

Formulae (3.14) and (3.18), which refer to the asymmetries in the standard model at $\sqrt{s} = m_Z$ in the narrow width approximation, can be derived from (3.28) and (3.29) by neglecting the V^0 contribution.

In all quantities, the only radiative correction included is the renormalization of α . For instance, in comparing the two models with $\Gamma(Z \rightarrow \ell^+\ell^-)$, δ_e in (3.8) has been neglected. In the asymmetries, δA^{weak} has been neglected, as well as the effect of real and virtual photon emission, given by δA^{QED} . On the other hand, in A^{Born} , the radiatively-corrected formulae for $\sin^2\theta$, (3.9) and (3.24)-(3.25) have been used.

4. - RESTRICTIONS ON THE PARAMETER SPACE

We now pass to the study of the expected restrictions on the parameter space of the model. We have considered the following quantities [16]:

i)	$\left(\frac{m_W}{m_Z} \right)$	at $m_Z = 92.1 \text{ GeV}$
ii)	$\Gamma(Z \rightarrow \ell^+\ell^-)$	at $m_Z = 92.1 \text{ GeV}$
iii)	$A_{\text{FB}}^{e^+e^- \rightarrow \tau^+\tau^-}, P_e = 0$	$\left\{ \begin{array}{ll} \text{at } \sqrt{s} = 92.1 \text{ GeV} & \text{LEP I} \\ \text{at } \sqrt{s} = 190 \text{ GeV} & \text{LEP II} \end{array} \right.$

$$\text{iv) } A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-} \quad p_e = 0.5 \quad \text{at } \sqrt{s} = 92.1 \text{ GeV}$$

$$\text{v) } A_{LR}^{e^+e^- \rightarrow \mu^+\mu^-} \quad \text{at } \sqrt{s} = 92.1 \text{ GeV}$$

We shall make the working assumption that these quantities are measured with the following precision:

$$\delta(m_W | m_Z) = \pm 2 \cdot 10^{-3} \quad (4.1)$$

$$\frac{\delta\Gamma(Z \rightarrow l^+l^-)}{\Gamma(Z \rightarrow l^+l^-)} = 0.02 \quad (4.2)$$

$$\delta A_{FB}(m_Z) = 0.01 \quad (4.3)$$

$$\delta A_{FB}(190 \text{ GeV}) = 0.03 \quad (4.4)$$

$$\delta A_{LR}(m_Z) = 0.02 \quad (4.5)$$

On such a basis we shall examine the corresponding restrictions on the model. We must comment that only (4.1), (4.2) and (4.4) reflect our previous analysis. In (4.3) and (4.5), the "theoretical uncertainties" given by the terms δA^{weak} have been superimposed on the experimental expectations (3.13) and (3.17). (δA^{QED} is detector-dependent and one hopes that the experimental data will be given with δA^{QED} already subtracted.) Our way of presenting the analysis in this section will be to assume that no deviation with respect to the standard model will be observed within the given precisions. In this case, (4.1) to (4.5) will give rise to limitations on the parameter space of the model, which is three-dimensional ($m_V, b, g/g''$).

We shall illustrate the possible limitations expected from different e^+e^- tests, choosing in each case the most convenient presentation. In Fig. 1, we present plots of m_W/m_Z versus m_Z for $m_V = 250 \text{ GeV}$ and three values of g/g'' : 0.05, 0.1 and 0.15. For comparison, we have also included the SM prediction. We observe that m_W/m_Z is always smaller than in the SM and decreases with increasing values of g/g'' . In the limit $g/g'' \rightarrow 0$, we recover the SM prediction. Assuming that no differences from the SM appear in the m_W/m_Z measurement, we can convert Fig. 1 into contours limiting the allowed range in the $(m_V, g/g'')$ plane. The curves in Fig. 2 represent the upper bound for g/g'' . For the error given in Eq. (4.1), we find $g/g'' \lesssim 0.07$. For a more pessimistic precision of $\delta(m_W/m_Z) = 10^{-2}$, we obtain

$g/g'' \lesssim 0.16$. In Fig. 3, the same test provides a strong restriction in a plane $(b, g/g'')$ at $m_V = 250$ GeV. The admissible region is given by a strip situated in the positive b range.

In Figs. 4 and 5 we plot the leptonic width $\Gamma(Z \rightarrow \ell^+\ell^-)$ versus m_Z for $m_V = 250$ GeV, $g/g'' = (0.1, 0.2, 0.3)$ and $b = 0, b = -0.05$ respectively. Some of the curves intersect the SM line and this circumstance allows for both small and large deviations from the SM, according to the actual value of m_Z . The limit $b \rightarrow 0, g/g'' \rightarrow 0$ gives the SM prediction. Here again, assuming that no discrepancies with the SM are found in measurements, one will obtain limits in the parameter space. Figure 6 shows the allowed region in a $(b, g/g'')$ plane for $m_V = 250$ GeV, $m_Z = 92.1$ GeV and for the precision given in Eq. (4.2). When g/g'' is sufficiently small, $\Gamma(Z \rightarrow \ell^+\ell^-)$ is particularly sensitive to b . For instance, at $g/g'' = 0.04$, we find $-0.025 < b < 0.04$. No severe limitations are found on a $(m_V, g/g'')$ plane at $b = 0$.

The asymmetries are plotted versus $\sqrt{s}$ in Figs. 7, 8 and 9 for an ideal acceptance ($x = 1$) and an illustrative choice of parameters: $m_V = 250$ GeV, $b = -0.04$, $g/g'' = 0.02$. Figure 7 refers to unpolarized ($P_e = 0$) A_{FB} . Figure 8 illustrates the effect of a 50% polarized electron beam on A_{FB} and Fig. 9 shows A_{LR}/P_e . SM lines are also displayed for comparison. Serious limitations on the parameter space can come from the measurement of the asymmetries if the predictions of the SM are verified within the aimed-for experimental precisions given in Eqs. (4.3)-(4.5). The forward-backward asymmetry with an unpolarized electron beam ($P_e = 0$) at $\sqrt{s} = m_Z$ provides the contoured region of Figs. 10 and 11. In Fig. 10, on a plane $(m_V, g/g'')$ at $b = 0$, the allowed region consists of two strips; the lower one requires $g/g'' < 0.12$. In Fig. 11, at $m_V = 250$ GeV in a plane $(b, g/g'')$, a more elaborate contour provides a strong constraint on the admitted region. These constraints become even more stringent with the use of a 50% longitudinally-polarized electron beam, as is shown in Fig. 12. Here, in a $m_V = 250$ GeV $(b, g/g'')$ plane, one can see that negative values of b are almost completely excluded. On a $(m_V, g/g'')$ plane we find $g/g'' < 0.08$ at $b = 0$. Less severe limitations are found from the measurement of A_{FB} at $\sqrt{s} = 190$ GeV. Here, in a $(m_V, g/g'')$ plane at $b = 0$, no additional constraints are found with respect to those previously discussed for $\sqrt{s} = m_Z$, apart from the case where m_V is close to 190 GeV and g/g'' is small, which gives rise to appreciable deviations from the SM predictions (an unpolarized electron beam is considered, but things do not really change, even assuming that polarization at that energy is available). Finally, in Fig. 13 the physical region selected by measurement of A_{LR} at $\sqrt{s} = m_Z$ with $P_e = 0.5$ is displayed. Figure 13 again provides, at $m_V = 250$ GeV, a strong limitation on b : $b > -0.01$. Moreover, we find $g/g'' \leq 0.12$ at $b = 0$.

In conclusion, one can say that the strongest constraints on g/g'' are expected to come from the measurements of m_W/m_Z with a precision of $\delta(m_W/m_Z) \equiv 2 \cdot 10^{-3}$ and $A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}$ at $\sqrt{s} = m_Z$ with $P_e = 0.5$ and $\delta(A_{FB}) \sim 0.01$: g/g'' must be smaller or equal to 0.08 if no deviations are observed. On the other hand, b will be strongly limited by the measurements of m_W/m_Z [assuming $\delta(m_W/m_Z) = 2 \cdot 10^{-3}$], $A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}$ at $\sqrt{s} = m_Z$ with $P_e = 0.5$ [assuming $\delta(A_{FB}) \sim 0.01$] and $A_{LR}^{e^+e^- \rightarrow \mu^+\mu^-}$ at $\sqrt{s} = m_Z$ with $P_e = 0.5$ [assuming $\delta(A_{LR}) \sim 0.02$]: a negative b can be almost completely ruled out if no deviations appear.

5. - SIGNATURES AT HADRON COLLIDERS

The UA1 experiment has recently quoted [6] the existence of two events in their $W^+ \rightarrow e^+\nu$ and $W^+ \rightarrow \mu^+\nu$ samples at very large transverse momentum ($p_{\perp}^W \approx 100$ GeV), in which the W^+ recoils against an energetic dijet system whose mass is compatible with that of a W or a Z ($70 \leq M_{jj} \leq 100$ GeV) and $250 \leq M_{\ell^+\nu jj} \leq 320$ GeV. These events are naturally expected in the SM from QCD double bremsstrahlung in weak boson production. However, the number of events predicted is 0.05 ± 0.03 , i.e., too small by two orders of magnitude to explain the UA1 sample. Heavy flavour production with $M_Q > M_W$ can also produce such events. Two recent evaluations [17,18] have shown that the predicted event rate is similar to QCD double bremsstrahlung expectations: for $M_Q \sim 90$ GeV/c, the expected number of events is 0.02. Since the mass of the dijet system corresponds to the mass of W or Z within two σ , the events can be due to standard WW , WZ or WH production (with $m_H \approx m_W$). Unfortunately, the expected cross-section is equal to 0.02 pb. Although we are aware that these events can be due to a non-Gaussian fluctuation in the response of the UA1 calorimeter or a statistical fluctuation, they might indicate some evidence beyond the standard model. It has been conjectured [19] that these events are mediated by a new particle having a mass of the order of 250 GeV and a width of 25 GeV which decays predominantly into W pairs. The new Z_3 arising from E_6 superstring-inspired models is a natural candidate. However, the predicted rate, as one can calculate from Ref. [20], is of the order of $4 \cdot 10^{-3}$ pb for $\ell\nu jj$ events, since the Z_3 couples to W pairs through mixing with Z_0 . It has recently been shown [21] that the new vector resonance V expected from the strongly-interacting Higgs sector of the SM can give the observed rate, provided one makes a convenient choice of the parameters.

The natural tool for detecting heavy bosons in hadronic collisions is the Drell-Yan mechanism which has been shown [22] to be very efficient in the discovery of the W and Z gauge bosons, and whose predictions from perturbative QCD are in agreement with UA1 and UA2 observations. Therefore the total cross-section for V

production in $p\bar{p}$ or pp collisions, which is obtained from convolution of the structure functions of the incoming hadrons with the partonic cross-section $\hat{\sigma}$ ($q\bar{q} \rightarrow V$) reads:

$$\sigma(H_1 H_2 \rightarrow V) = \frac{4\pi^2 \alpha}{3s} K(m_V) \int_{\frac{m_V^2}{s}}^1 \frac{dx}{x} (g_V^{q^2} + g_A^{q^2}) \left[q^{H_1}(x, m_V^2) \bar{q}^{H_2}\left(\frac{m_V^2}{xs}, m_V^2\right) + (q \leftrightarrow \bar{q}) \right] \quad (5.1)$$

where K is the factor due to QCD corrections and $eg_V^q = -\frac{1}{2}v_2^q$, $eg_A^q = \frac{1}{2}a_2^q$, with v_2^q and a_2^q given in (2.39) (for the neutral vector boson) and $eg_V^q = -eg_A^q = (1/2\sqrt{2})h_V^q$, with h_V^q given in (2.31) (for the charged one). We will use for K the approximate expression obtained from soft gluon resummation [23]:

$$K(m_V) \simeq \exp \left[\frac{\alpha_s(m_V^2)}{2\pi} \frac{4}{3} \pi^2 \right] \quad (5.2)$$

For the quark structure functions we have used the parametrization of Duke and Owens [24]. Since predictions for lepton pair production at the Tevatron and SSC have been previously computed [25], we will concentrate here on W and Z pair production. This cross-section is directly obtained from Eq. (5.1) by multiplying $\sigma(H_1 H_2 \rightarrow V)$ with the appropriate branching ratios. For W^+W^- production from V^0 , we get:

$$\Gamma_{V^0 \rightarrow W^+W^-} = \frac{m_{V^0}}{12\pi} \frac{g_{V^0W^+W^-}^2}{16} \left(1 - 4 \frac{m_W^2}{m_{V^0}^2} \right)^{3/2} \left(\frac{m_{V^0}}{m_W} \right)^4 \left[1 + 20 \left(\frac{m_W}{m_{V^0}} \right)^2 + 12 \left(\frac{m_W}{m_{V^0}} \right)^4 \right] \quad (5.3)$$

and for WZ production from charged V :

$$\Gamma_{V^\pm \rightarrow W^\pm Z} = \frac{m_{V^\pm}}{12\pi} \frac{g_{V^\pm W^\pm Z}^2}{16} \left[\left(1 + \frac{m_W^2 - m_Z^2}{m_{V^\pm}^2} \right)^2 - 4 \frac{m_W^2}{m_{V^\pm}^2} \right]^{3/2} \frac{m_{V^\pm}^4}{m_Z^2 m_W^2} \left[1 + 10 \left(\frac{m_W^2 + m_Z^2}{m_{V^\pm}^2} \right) + \frac{m_W^4 + m_Z^4 + 10 m_W^2 m_Z^2}{m_{V^\pm}^4} \right] \quad (5.4)$$

For large values of g''/g , we will consider the coupling constants $g_{V^0 W^+W^-}$ and

g_{VWZ}^{++} to be dominated by higher derivative terms (see Section 2 for more details). It has been recently shown [21] that the anomalous UAl events, in which a W is produced at very large transverse momentum ($p_{\perp}^W \approx 100$ GeV) with a dijet system whose mass is near the W or Z mass, are compatible with a V of mass around 250 GeV decaying predominantly into W and Z pairs provided b is negative ($b \approx -0.05$) and g/g'' small (~ 0.02). Including radiative corrections and asking for a total cross-section in l^+vjj from V^0 production greater than 0.5 pb with $\Gamma_V \leq 25$ GeV, we plot in Fig. 14 the allowed region in the $(b, g/g'')$ plane for different values of γ , fixing m_V to 250 GeV. We note that the combination with low-energy constraints gives very strong limitations to the allowed $(b, g/g'')$ domain. A striking check for the consistency of the model, if UAl events are confirmed, will be provided by measurements of the ratio m_W/m_Z , the forward-backward asymmetry A_{FB} and left-right asymmetry A_{LR} at LEP and SLC, since in this parameter domain we expect strong deviations from the standard model predictions as shown in Figs. 3, 12 and 13. For the asymmetries, this is due to the fact that for negative b values, the couplings of the physical vector bosons W and Z to fermion pairs differ significantly from the standard ones.

We will now discuss the M_{Wjj} distribution:

$$\frac{d\sigma}{dM_{Wjj}} = \frac{8\pi\alpha}{3s} K(M_{Wjj}) \frac{M_{Wjj}^3}{m_V} \frac{\Gamma_{V \rightarrow W^+W^-, ZW}}{(M_{Wjj}^2 - m_V^2)^2 + m_V^2 \Gamma_V^2}$$

$$\int_{\frac{M_{Wjj}}{s}}^1 \frac{dx}{x} \left[q^{H_1}(x, M_{Wjj}^2) \bar{q}^{H_2}\left(\frac{M_{Wjj}^2}{xs}, M_{Wjj}^2\right) + (q \leftrightarrow \bar{q}) \right] (g_V^q{}^2 + g_A^q{}^2) \quad (5.5)$$

As shown in Fig. 15 for $m_V = 250$ GeV, the distributions drop by roughly a factor of 50 between $M_{Wjj} = 250$ GeV and $M_{Wjj} = 300$ GeV at SppS collider energy and roughly by a factor of 20 at Tevatron energy ($\sqrt{s} = 1.8$ TeV). We predict twenty times more events at the Tevatron collider than at the CERN collider.

We will next consider the transverse momentum distribution of the W^+ . In the Born approximation, the differential cross-section reads [26]:

$$\left. \frac{d\sigma^{\text{BORN}}}{dP_{\perp}^W dy} \right|_Y = \frac{2P_{\perp}^W}{3} \int_{x_{\min}}^1 dx_a \frac{1}{x_a - x_{\perp} \left(\frac{\chi + \cos \theta}{2 \sin \theta} \right)} x_a x_b$$

$$\left[q^{H_1}(x_a, m_V^2) \bar{q}^{H_2}(x_b, m_V^2) + (q \leftrightarrow \bar{q}) \right] \frac{d\hat{\sigma}}{dt}$$

(5.6)

with

$$x_{\perp} = \frac{2P_{\perp}^W}{\sqrt{s}}, \quad \chi = \sqrt{1 + \frac{4m_W^2 \sin^2 \theta}{x_{\perp}^2 s}}$$

$$\sin \theta = \frac{P_{\perp}^W}{\sqrt{P_{\perp}^2 \cosh^2 y + m_W^2 \sinh^2 y}}, \quad x_b = \frac{2\Delta + x_a x_{\perp} (\chi - \cos \theta) / \sin \theta}{2x_a - x_{\perp} (\chi + \cos \theta) / \sin \theta}$$

$$x_{\min} = \frac{x_{\perp} (\chi + \cos \theta) / \sin \theta + 2\Delta}{2 - x_{\perp} (\chi - \cos \theta) / \sin \theta}$$

For W^+ production from V^0 we get:

$$\frac{d\hat{\sigma}}{dt} = 2\pi\alpha^2 \left(\frac{g_{V^0 W^+ W^-}^2}{e^2} \right) (g_V^2 + g_A^2) \frac{1}{(\hat{s} - m_V^2)^2 + m_V^2 \Gamma_V^2}$$

$$\left\{ \left(\frac{\hat{u} \hat{t}}{m_W^4} - 1 \right) \left[\frac{1}{4} - \frac{m_W^2}{\hat{s}} + 3 \frac{m_W^4}{\hat{s}^2} \right] + \frac{\hat{s}}{m_W^2} - 4 \right\}$$

(5.7)

with $\Delta = 0$, whereas for W^+ production from V^+ , we get:

$$\begin{aligned} \frac{d\hat{\sigma}}{dt} = & 2\pi\alpha^2 \left(\frac{g_{V+W+Z}^2}{e^2} \right) (g_V^2 + g_A^2) \frac{1}{(\hat{s} - m_V^2)^2 + m_V^2 \Gamma_V^2} \\ & \left\{ \left(\frac{\hat{u}\hat{t}}{m_W^2 m_Z^2} - 1 \right) \left[\frac{1}{4} - \frac{m_W^2 + m_Z^2}{2\hat{s}} + \frac{(m_W^2 + m_Z^2)^2 + 8m_W^2 m_Z^2}{4\hat{s}^2} \right] \right. \\ & \left. + \frac{m_W^2 + m_Z^2}{m_W^2 m_Z^2} \left[\frac{\hat{s}}{2} - m_W^2 - m_Z^2 + \frac{(m_W^2 - m_Z^2)^2}{2\hat{s}} \right] \right\} \end{aligned} \quad (5.8)$$

with $\Delta = (m_Z^2 - m_W^2)/s$. In Eqs. (5.7) and (5.8) the Mandelstam variables $\hat{s}$, $\hat{t}$, $\hat{u}$ read:

$$\begin{aligned} \hat{s} &= x_a x_b s \\ \hat{t} &= m_W^2 - x_a x_{1\perp} s \left(\frac{x - \cos\theta}{2\sin\theta} \right) \\ \hat{u} &= m_W^2 - x_b x_{1\perp} s \left(\frac{x + \cos\theta}{2\sin\theta} \right) \end{aligned}$$

As shown in Fig. 16, for $\sqrt{s} = 630$ GeV, $(d\sigma^{\text{Born}}/dp_{\perp}^W dy)|_{y=0}$ is maximal around $p_{\perp}^W \approx 90$ GeV and exceeds the QCD prediction of Altarelli et al. [27] by roughly a factor of seven in this large $p_{\perp}^W$ range. Similar conclusions can be drawn for W^+ production from a V^+ , which gives an excess of events of roughly a factor 3. We now have to take into account the effect of soft gluon radiation from quark and antiquark legs. We will use the approximate formula [28]

$$\left. \frac{d\sigma}{dp_{\perp}^W dy} \right|_y = \int_0^{p_{\perp}^W} dk_{\perp} \left. \frac{d\sigma^{\text{Born}}}{dk_{\perp} dy} \right|_y \left. \frac{dP}{dk} \right|_{k=p_{\perp}^W - k_{\perp}} \quad (5.9)$$

where the soft gluon transverse momentum distribution is given by [29]:

$$\frac{dP}{dk^2} = \frac{1}{2} \int_0^{\infty} b db J_0(bk) \exp \Delta(b) \quad (5.10)$$

with

$$\Delta(b) = \frac{8}{3\pi} \int_0^{m_V e^{-3/2}} \frac{dq}{q} \left(J_0(bq) - 1 \right) d_s(q) \left(\ln \frac{m_V^2}{q^2} - \frac{3}{2} \right) \quad (5.11)$$

As shown in Fig. 16 (full curves), the effect of soft gluon radiation is to shift the Jacobian peak by a few GeV and to increase the cross-section in the large $p_{\perp}^W$ range ($p_{\perp}^W \gtrsim 100$ GeV).

Let us remark that we have not attempted a direct comparison with the experimental quantity $(1/N)(dN/dp_{\perp}^W)$ [6], although acceptance and resolution are corrected, for two reasons. First, the theoretical cross-section has been computed at a fixed rapidity and the absolute normalization depends strongly on structure functions which are non-perturbative ingredients. Secondly, we are interested in the ratio between the QCD background and the excess of events in the large $p_{\perp}^W$ domain, which is quite independent of such uncertainties. We can therefore conclude that our predictions are in agreement with UA1 data. We will now give predictions for the Tevatron collider, taking for the parameters b , g/g'' and γ the values needed to explain UA1 anomalous events with $m_V = 250$ GeV and $\Gamma_V < 25$ GeV. As shown in Fig. 17 the signal exceeds the QCD background for $p_{\perp}^W > 85$ GeV. If UA1 events are not a statistical fluctuation, they should be seen at the Tevatron. If we decrease the V mass (let us take for example $m_{V^0} = 200$ GeV), the $p_{\perp}^W$ is shifted to lower transverse momentum values. As shown in Fig. 18, at $\sqrt{s} = 630$ GeV the peak is around 60 GeV, its magnitude being dependent on the b and g/g'' values chosen. Taking $b = 0$ and $g/g'' = 0.08$, our predictions become closer to expectations from $t\bar{t}$ production with a top mass just above the W mass [17,18].

Let us assume now that these UA1 anomalous events are just a statistical fluctuation and will disappear with ACOL luminosity. If no deviation from SM expectations for the ratio m_W/m_Z and for the asymmetries A_{FB} and A_{LR} is seen at LEP/SLC, this will indicate that $b > 0$ and $g/g'' < 0.08$. We therefore have to look for V production at LHC and SSC. As an illustration, we will take $b = 0$ and $g/g'' = 0.05$.

We give in Fig. 19 the W jet-jet mass distribution for $m_{V^0} = 250$ GeV at LHC collider energy $\sqrt{s} = 17$ TeV, assuming two characteristic values for the V^0 width, $\Gamma_{V^0} = 10$ GeV and $\Gamma_{V^0} = 157$ GeV. As can be seen from inspection of this figure, the background due to electroweak WW production is important. It has been extracted from Ref. [30] at $\sqrt{s} = 20$ TeV. At the SSC collider, the WW signal from V^0 lies under the standard electroweak background. We should note that angular distributions should certainly help to distinguish between WW production from a V^0 and standard production, since in the former case the decay is isotropic, whereas in the latter case the angular distribution is peaked around $|\cos \theta| = 1$. We give in

Fig. 20 predictions for $m_{V^0} = 500$ GeV at the LHC collider. Similar curves are obtained for the SSC collider. In Fig. 21 (resp. 22) we give the W jet-jet mass distribution at SSC collider energy ($\sqrt{s} = 40$ TeV) for $m_{V^0} = 1$ TeV (resp. 2 TeV). Now the signal exceeds the standard electroweak background if the V^0 is not too wide. Similar conclusions can be obtained for the LHC collider.

Total cross-sections have the same magnitude as expected in E_6 superstring-inspired models [20]. However, our predictions are certainly underestimated at high energy, since we have not included the contribution due to WW fusion. This fusion mechanism is not expected to increase the standard electroweak background by more than 20% in the absence of Higgs production, since $W_T W_T$ production is the dominant process from quark-antiquark annihilation.

6. - CONCLUSIONS

We have discussed possible high-energy tests of an electroweak model which can be seen as an effective description of the SM in the limit when the parameter m_H , the Higgs mass of the electroweak Lagrangian, becomes very large. For such a limit, the longitudinal scalar sector of the SM becomes strong-interacting. The use of our electroweak model may also be seen as a way of parametrizing possible deviations from the SM in terms of bounds on a definite set of parameters having a possible physical significance. Also, taking a more radical attitude, nothing prevents us from considering the model as an alternative to the SM itself. The meaning of such an alternative can be expressed as follows: whereas the spontaneous electroweak symmetry breaking is realized linearly in the SM, in our model it is instead realized non-linearly. The main advantage of the non-linear realization is that since it is entirely expressed through the co-ordinates of the quotient space (initial symmetry group over surviving symmetry group), it has no explicit degrees of freedom left for the physical Higgs particle. So the physical Higgs particle disappears as an elementary field.

The disadvantage of the non-linear approach is that it allows only for an effective Lagrangian approach. One can easily see that, under general conditions, in a linear realization of a spontaneous symmetry breaking there is no way of choosing the representation for the scalar fields such that they all become absorbed to give masses to the gauge bosons. If this were possible, each direction away from the minimum of the potentials would have to correspond to a Goldstone mode implying a totally flat potential. The non-linear realization considered here is almost an exact replica of the QCD effective Lagrangian description via a non-linear σ -model. We expect bound states corresponding to the ρ -mesons to appear. The main difference is that the pions, which are the co-ordinates of the quotient

space, and which are physical particles (Goldstones) in QCD with massless quarks, are here absorbed to give masses to W and Z.

The picture when extended within the QCD paradigm leads, as is well known, to technicolour, which requires a more complex setting in order to be consistent. Here we expect that, leaving out any QCD analogy and contenting ourselves with the non-linear effective Lagrangian approach, this picture may correspond to the strong interacting bosonic sector. The model has no surviving Higgs degrees of freedom, but it has instead a triplet of massive vector bosons V^+ , V^0 , V^- , whose mass we are not able to predict theoretically, and which mix with the gauge bosons of $SU(2)_L \times U(1)$, W, Z, γ . The model being dynamically different from technicolour, one cannot apply the usual scaling arguments from QCD which leads to a large mass of the ρ of technicolour. Through mixing, V^+ , V^0 , V^- get coupled to the fermions. We have called g'' the gauge coupling of the V's. The standard model results are reproduced for $g'' \rightarrow \infty$. In addition to the couplings arising from the mixing, it is, however, possible to imagine a direct coupling of V to the fermions without violating the symmetry of the problem. In particular, the couplings to the left-handed fermions are described by a new parameter b which should, in principle, be calculable from V-W mixing and one-loop corrections to the W-fermion vertex and to the W propagator. From these graphs, b is expected to be small, of the order of $g^2(g/g'')^2$. A complete study of the low-energy phenomenology within the model has been performed and it has led to stringent bounds on the parameters of the model [5], corresponding to the agreement of the standard model with observation.

In this paper, we have first discussed the possible incidence of precise e^+e^- measurements from SLC and LEP. We have concluded that measurements of m_W/m_Z with a precision $\delta(m_W/m_Z) \cong 2 \cdot 10^{-3}$ will give the strongest constraints on the ratio g/g'' , together with accurate (one per cent) measurements of the forward-backward asymmetry in $e^+e^- \rightarrow \mu^+\mu^-$ at the Z mass with longitudinally-polarized electrons (degree of polarization $\sim \frac{1}{2}$). If no deviations are observed within these limits, one will have g/g'' smaller or equal to eight per cent. The parameter b will be strongly limited by the above-mentioned measurements of m_W and m_Z , by one per cent measurements of the forward-backward asymmetry in $e^+e^- \rightarrow \mu^+\mu^-$ at the Z mass, again with 50 per cent longitudinally-polarized electrons, and by measurements of the left-right asymmetry in the same reaction, always at the Z mass, and with similarly polarized electrons if a two per cent accuracy can be reached. In particular, a negative value of b would be almost excluded for no deviations. A more complete comparison could be made by considering the whole set of forward-backward asymmetries in $e^+e^- \rightarrow \mu^+\mu^-$, $\tau^+\tau^-$, $c\bar{c}$, $b\bar{b}$, etc. at the Z^0 from e^+e^- beams, especially if longitudinally polarized. The comparative study of these asymmetries would allow us uniquely to disentangle effects which are typical of this model from effects which are foreseen in the standard model and to overdetermine the parameters involved.

Concerning proton colliders, we have seen that the interpretation of $\ell\nu jj$ anomalous events from a V resonance of the strongly-interacting Higgs sector of the SM requires a direct coupling of the V to the fermions which has to be negative. If these events are not a statistical fluctuation, they need some confirmation at the CERN collider with ACOL luminosity and at the Tevatron collider. Similar $\ell\nu jj$ events are expected in the W^- sample with the same rate. We have shown that a further crucial test for the theory will be provided by measurements of the ratio m_W/m_Z and of the forward-backward and left-right asymmetries at the Z peak at LEP/SLC, since if the anomalous events are confirmed and if our interpretation of them is correct, we expect strong deviations from the SM predictions. If these anomalous events do not completely disappear, but shift to lower W transverse momentum values, they can be qualitatively explained by W pair production from a V resonance if the V mass is smaller. However, this mechanism will compete with W pair production from $t\bar{t}$ if the top mass is higher than the W one. Nevertheless, the topology of these events will be different since the leptonic decay of the W produced from a top quark is always accompanied by a hadronic jet coming from the b quark. If we are pessimistic, i.e., we assume that these events will disappear with ACOL luminosity, and that nothing will be seen at the Tevatron and that W and Z masses and asymmetries at LEP/SLC will be in agreement with SM expectations, then, as we have shown, detection of the V boson through its decay into W pairs is possible at LHC and SSC. If the V is heavy ($m_V > 500$ GeV), and not too wide, it can be directly seen from the WW mass distribution. The background due to standard electroweak WW production will be dominant for lower V masses if the width of the V is large. Nevertheless, angular distributions should certainly help to discriminate between the two mechanisms.

ACKNOWLEDGEMENTS

We would like to thank D. Denegri, J. Ellis, M. Greco and G. Martinelli for interesting and stimulating discussions, and D. Cocolicchio for his technical advice on the computer graphics. This work was partially supported by the Swiss National Science Foundation.

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FIGURE CAPTIONS

- Fig. 1 : The ratio m_W/m_Z versus m_Z in the presence of the vector bosons V (solid lines) for $m_V = 250$ GeV and:
- a) $g/g'' = 0.05$
 - b) $g/g'' = 0.1$
 - c) $g/g'' = 0.15$
- The dashed line represents the SM prediction.
- Fig. 2 : Contours limiting the allowed $(m_V, g/g'')$ range at $b = 0$ resulting from a measurement of m_W/m_Z with a precision of:
- a) 10^{-2}
 - b) $2 \cdot 10^{-3}$
- The shaded areas are the allowed ones.
- Fig. 3 : Allowed region (shaded area) in the $(b, g/g'')$ plane at $m_V = 250$ GeV, coming from a measurement of m_W/m_Z with a precision of:
- a) 10^{-2}
 - b) $2 \cdot 10^{-3}$

Fig. 4 : The leptonic width of Z, $\Gamma(Z \rightarrow \ell^+\ell^-)$ versus m_Z in the presence of the vector bosons V (solid lines) for $m_V = 250$ GeV, $b = 0$ and:

- a) $g/g'' = 0.1$
- b) $g/g'' = 0.2$
- c) $g/g'' = 0.3$

The dashed line represents the SM prediction.

Fig. 5 : The leptonic width of Z, $\Gamma(Z \rightarrow \ell^+\ell^-)$ versus m_Z in the presence of the vector boson V (solid lines) for $m_V = 250$ GeV, $b = -0.05$ and:

- a) $g/g'' = 0.1$
- b) $g/g'' = 0.2$
- c) $g/g'' = 0.3$

The dashed line represents the SM prediction.

Fig. 6 : Allowed region (shaded area) in the $(b, g/g'')$ plane at $m_V = 250$ GeV coming from a measurement of $\Gamma(Z \rightarrow \ell^+\ell^-)$ with a 2% precision. The Z mass was chosen to be 92.1 GeV.

Fig. 7 : The forward-backward asymmetry $A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}$ for an unpolarized electron beam ($P_e = 0$) and ideal detector acceptance ($x = 1$) for $m_V = 250$ GeV, $b = -0.04$ and $g/g'' = 0.02$. The dashed line represents the SM prediction. The Z mass was chosen to be 92.1 GeV.

Fig. 8 : The forward-backward asymmetry $A_{FB}^{e^+e^- \rightarrow \mu^+\mu^-}$ for ideal detector acceptance ($x = 1$), $P_e = 0.5$, $m_V = 250$ GeV, $b = -0.04$ and $g/g'' = 0.02$. The dashed line represents the SM prediction. For the Z mass we used 92.1 GeV.

Fig. 9 : The left-right asymmetry $A_{LR}^{e^+e^- \rightarrow \mu^+\mu^-}/P_e$ for $m_V = 250$ GeV, $b = -0.04$ and $g/g'' = 0.02$. The dashed line represents the SM prediction. For the Z mass we used 92.1 GeV.

Fig. 10 : Allowed region (shaded area) in the $(m_V, g/g'')$ plane at $b = 0$ coming from a measurement of $A_{FB}(\sqrt{s} = m_Z)$ with an error $\delta A_{FB} = 0.01$ and an unpolarized electron beam ($P_e = 0$). We used $m_Z = 92.1$ GeV.

Fig. 11 : Allowed region (shaded area) in the $(b, g/g'')$ plane at $m_V = 250$ GeV resulting from a measurement of $A_{FB}(\sqrt{s} = m_Z)$ with an error $\delta A_{FB} = 0.01$ and an unpolarized electron beam ($P_e = 0$). We used $m_Z = 92.1$ GeV.

Fig. 12 : Allowed region (shaded area) in the $(b, g/g'')$ plane at $m_V = 250$ GeV resulting from a measurement of $A_{FB}(\sqrt{s} = m_Z)$ with an error $\delta A_{FB} = 0.01$ and a polarized electron beam ($P_e = 0.5$). We used $m_Z = 92.1$ GeV.

Fig. 13 : Allowed region (shaded area) in the $(b, g/g'')$ plane at $m_V = 250$ GeV resulting from a measurement of A_{LR} ($\sqrt{s} = m_Z$) with an error $\delta A_{LR} = 0.02$ and a 50% longitudinally-polarized electron beam ($p_e = 0.5$). We used $m_Z = 92.1$ GeV.

Fig. 14 : Allowed region for $m_{V^0} = 250$ GeV, $\Gamma_{V^0} < 25$ GeV and $\sigma_{\ell^+\nu jj} > 0.5$ pb in the $(b, g/g'')$ plane:

- a) $\gamma = 0.125$
- b) $\gamma = 0.1$
- c) $\gamma = 0.075$
- d) $\gamma = 0.05$

The shaded area represents the region allowed by the low-energy tests [5].

Fig. 15 : The W jet-jet mass distribution at $\sqrt{s} = 630$ GeV and $\sqrt{s} = 1.8$ TeV from V^0 production (full curves) and V^+ production (dashed curves).

Fig. 16 : The transverse momentum distribution $d\sigma/dp_{\perp}^W dy$ at $y = 0$ for the Sp $\bar{p}$ S collider. Thick curve: QCD contribution from Ref. [27]. Dashed curves: Born contribution. Full curves: soft gluon radiation included. The decay $W^+ \rightarrow \ell^+ \nu$ has been included.

Fig. 17 : Same as Fig. 4 for the Tevatron collider ($\sqrt{s} = 1.8$ GeV).

Fig. 18 : Transverse momentum distribution $(d\sigma/dp_{\perp}^W dy)|_{y=0}$ for $m_{V^0} = 200$ GeV, $\Gamma_{V^0} = 25$ GeV, $b = 0$, $g/g'' = 0.08$ at $\sqrt{s} = 630$ GeV. Thick curve: QCD contribution from Ref. [27]. Full curve: predictions from V^0 production with soft gluon radiation included.

Fig. 19 : The W jet-jet mass distribution for $m_{V^0} = 250$ GeV at the LHC collider ($\sqrt{s} = 17$ TeV) with $b = 0$ and $g/g'' = 0.05$. Thick curve: standard electroweak contribution from $q\bar{q}$ at $\sqrt{s} = 20$ TeV (Ref. [30]). Full curve: $\Gamma_{V^0} = 10$ GeV. Dashed curve: $\Gamma_{V^0} = 157$ GeV.

Fig. 20 : The W jet-jet mass distribution for $m_{V^0} = 500$ GeV at the LHC collider ($\sqrt{s} = 17$ TeV) for $b = 0$ and $g/g'' = 0.05$. Thick curve: standard electroweak contribution at $\sqrt{s} = 20$ TeV (Ref. [30]). Full curve: $\Gamma_{V^0} = 32$ GeV. Dashed curve: $\Gamma_{V^0} = 130$ GeV.

Fig. 21 : The W jet-jet mass distribution for $m_{V^0} = 1$ TeV at the SSC collider ($\sqrt{s} = 40$ TeV). Thick curve: standard electroweak contribution at $\sqrt{s} = 40$ TeV (Ref. [30]). Full curve: $\Gamma_{V^0} = 34$ GeV. Dashed curve: $\Gamma_{V^0} = 307$ GeV.

Fig. 22 : The W jet-jet mass distribution for $m_{V^0} = 2$ TeV at the SSC collider. Thick curve: standard electroweak background. Full curve: $\Gamma_{V^0} = 16$ GeV. Dashed curve: $\Gamma_{V^0} = 260$ GeV.

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