In summary, the decay $\tau^- \to K^- \nu_\tau$ has been observed at the level expected if the τ couples to the standard Cabibbo-suppressed axial-vector hadronic weak current.

This work was supported in part by the U. S. Department of Energy under Contracts No. DE-AC03-76SF00515 and No. W-7405-ENG-48. Support for individuals came from the listed institutions.

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[1]All reactions in this paper also imply the charge-conjugate reaction. Thus, $\tau^- \to K^- \nu_\tau$ stands for itself and for $\tau^+ \to K^+ \bar{\nu}_\tau$.

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Bounds on $M_{Z_{1,2}}$ in Any $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ Gauge Model

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(Received 29 March 1982)

If the known fermions transform under $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, then the two neutral vector gauge bosons must satisfy the bounds 83 GeV $< M_{Z_1} <$ 116 GeV and $M_{Z_2} >$ 200 GeV. In addition, if mixing is negligible for the two charged vector gauge bosons, then 75 GeV $< M_{W_1} <$ 97 GeV. These results come directly from analyzing existing experimental data and do not depend on any further theoretical assumptions, such as the value of g_L/g_R or the choice of Higgs-boson representations.

PACS numbers: 14.80.Er, 12.10.Ck, 12.30.Ez

In the standard $SU(2) \otimes U(1)$ electroweak gauge model,[1] the $W^\pm$ and Z^0 vector bosons are the mediators of the charged-current and neutral-current weak interactions, respectively. Their predicted masses are[2] $M_W = 83.0 \pm 2.4$ GeV and $M_Z = 93.18 \pm 2.0$ GeV, where radiative corrections to the lowest-order mass formulas $G_F/\sqrt{2} = e^2/(8M_W^2 \sin^2\theta_W)$ and $M_Z = M_W/\cos\theta_W$ have been taken into account. The new $p\bar{p}$ collider at CERN is expected to be capable of producing these particles at an observable rate in the near future. On the other hand, if the electroweak gauge group is really $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$,[3-14] then the $W^\pm$ and Z^0 bosons observable at about 100 GeV may not have the masses given above. It is therefore important to determine the allowed mass ranges for these bosons with as little extra theoretical input as possible. Accordingly, in the following analysis, we let the $SU(2)_L$ and $SU(2)_R$ couplings, g_L and g_R, be free parameters and we do not restrict ourselves to any particular set of Higgs bosons for the spontaneous symmetry breaking. Our results are summarized in Figs. 1 and 2.

Consider the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ with the electromagnetic current given by $J_{em} = J_{3L} + J_{3R} + \frac{1}{2}J_{B-L}$. The fundamental fermions are of course the quarks and leptons with baryon number B and lepton number L equal to

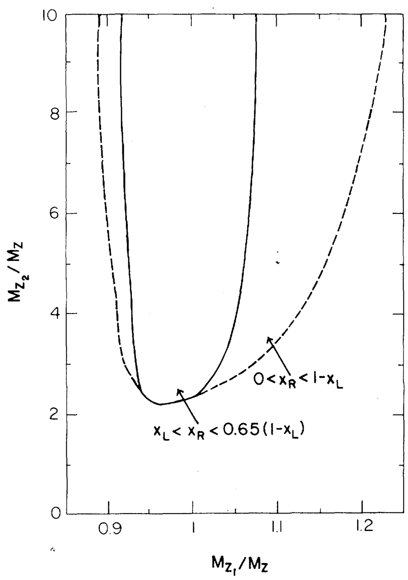

FIG. 1. Allowed regions of M_{Z_1} and M_{Z_2}, normalized to $M_Z = 93.8$ GeV. For the special case $x_L = x_R$, the bounds are $0.94 < M_{Z_1}/M_Z < 1.06$ and $M_{Z_2}/M_Z > 3.2$.

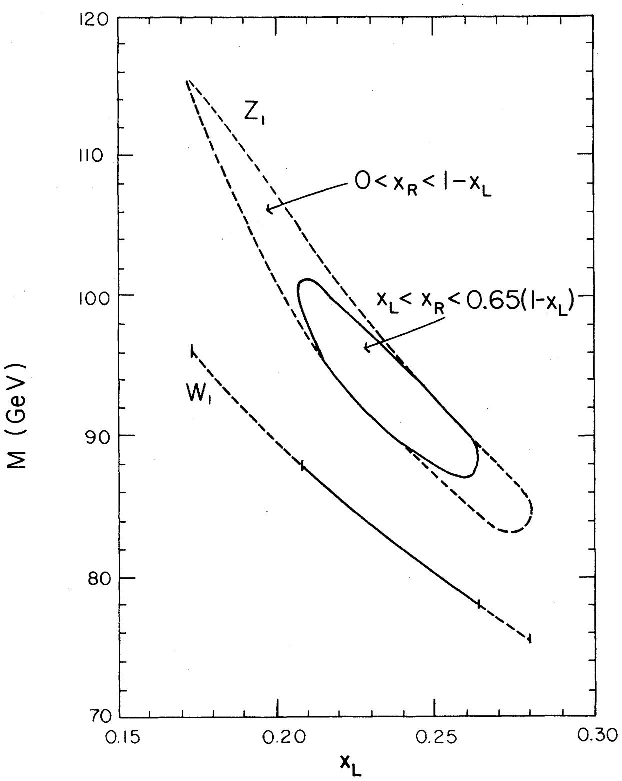

FIG. 2. Allowed Z_1 and W_1 masses vs the weak parameter x_L. For the special case $x_L = x_R$, M_{Z_1} is between 88 and 100 GeV. The allowed range of x_L is unchanged.

$(\frac{1}{3},0)$ and $(0,1)$, respectively. The left-handed (right-handed) fermions are doublets (singlets) under $SU(2)_L$ and singlets (doublets) under $SU(2)_R$. Let the neutral-current interaction be

$$H_{NC} = g_L W_{3L} J_{3L} + g_R W_{3R} J_{3R} + \tfrac{1}{2} g_C C J_{B-L}, \tag{1}$$

where g_L, g_R, g_C are gauge couplings and W_{3L}, W_{3R}, C are the corresponding vector gauge bosons. We do not assume any specific energy scale for the breaking of $SU(2)_R$ to $U(1)$, and hence g_R in Eq. (1) is defined only for the neutral-current sector and may be different from the corresponding coupling in the charged-current sector. On the other hand, a single g_L suffices for both neutral-current and charged-current sectors as in the standard model. Since $e^{-2} = g_L^{-2} + g_R^{-2} + g_C^{-2}$, we find it convenient to define $x_L \equiv e^2/g_L^2$ and $x_R \equiv e^2/g_R^2$ so that $x_L + x_R = 1 - e^2/g_C^2 < 1$. The photon is then given by $A = x_L^{1/2} W_{3L} + x_R^{1/2} W_{3R} + (1 - x_L - x_R)^{1/2} C$. We now define the orthogonal weak gauge fields

$$\begin{aligned} Z &\equiv (1-x_L)^{1/2} W_{3L} - x_L^{1/2}(1-x_L)^{-1/2}[x_R^{1/2} W_{3R} + (1-x_L-x_R)^{1/2} C], \\ D &\equiv (1-x_L)^{-1/2}[(1-x_L-x_R)^{1/2} W_{3R} - x_R^{1/2} C]. \end{aligned} \tag{2}$$

In terms of this new basis, the weak-interaction part of Eq. (1) is then

$$H_{NC}^{\text{weak}} = e[x_L(1-x_L)]^{-1/2} Z (J_{3L} - x_L J_{em}) + g_D D J_D, \tag{3}$$

where

$$g_D J_D = e(1-x_L-x_R)^{1/2} x_R^{-1/2}(1-x_L)^{-1/2} J_{3R} - \tfrac{1}{2} e\, x_R^{1/2}(1-x_L)^{-1/2}(1-x_L-x_R)^{-1/2} J_{B-L}. \tag{4}$$

Since g_D and J_D always appear together in the combination $g_D J_D$, we can choose for convenience $g_D \equiv e[x_L(1-x_L)]^{-1/2}$; then using the form of J_{em}, we find

$$J_D = (x_L x_R)^{1/2}(1-x_L-x_R)^{-1/2} J_{ZL} + (1-x_L) x_L^{1/2} x_R^{-1/2}(1-x_L-x_R)^{-1/2} J_{ZR}, \tag{5}$$

where $J_{ZL} \equiv J_{3L} - x_L J_{em}$ and $J_{ZR} \equiv J_{3R} - x_R J_{em}$.

To obtain the effective interaction at low energies, we consider the most general mass-squared matrix for the Z and D vector bosons:

$$M^2 = \frac{2e^2}{x_L(1-x_L)}\begin{bmatrix} A & B \\ B & C \end{bmatrix}. \tag{6}$$

Then

$$H_{\rm NC}^{\rm eff} = (4G_F/\sqrt{2})\{(\rho_1 J_{ZL})^2 + (\rho_2 J_{ZL} + \eta J_{ZR})^2\}, \tag{7}$$

where

$$(8G_F/\sqrt{2})\rho_1^2 = (2A)^{-1}, \tag{8}$$

$$\rho_2/\rho_1 = [C/A - B^2/A^2]^{-1/2}\{(x_L x_R)^{1/2}(1-x_L-x_R)^{-1/2} - B/A\}, \tag{9}$$

$$\eta/\rho_1 = [C/A - B^2/A^2]^{-1/2}(1-x_L)x_L^{1/2}x_R^{-1/2}(1-x_L-x_R)^{-1/2}. \tag{10}$$

Hence, the most general low-energy neutral-current interaction in $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ is described by five parameters: x_L, x_R, ρ_1, ρ_2, and η. In the limit where ρ_2 and η go to 0 and ρ_1 goes to 1, we recover the predictions of the standard model with x_L as the only free parameter.

In previous investigations of left-right extensions of the standard model, it was either assumed that x_R is equal to x_L (left-right *symmetric* models[3-8]) or that x_R is not equal to x_L (left-right *asymmetric* models[9-14]) but nevertheless determinable from a grand-unification condition, such as $x_R = \frac{3}{5}(1-x_L)$. In addition, most authors chose a specific set of Higgs bosons for the symmetry breaking and in so doing, fixed or restricted the parameters ρ_1, ρ_2, and η. For example, one popular choice[9-11] was $\rho_1 = 1$ and $\rho_2 = 0$. In this work, we are not committed to any such theoretical restriction. We simply analyze the existing experimental data on neutral currents in the context of the given effective Hamiltonian, i.e., Eq. (7), and obtain empirical constraints on all five parameters x_L, x_R, ρ_1, ρ_2, and η independently. These in turn provide us with constraints on M_{Z_1} and M_{Z_2} as well as M_{W_1} if the mixing in the charged-boson sector is negligible.

We now confront the most general effective neutral-current weak-interaction Hamiltonian for $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, i.e., Eq. (7), with experimental data. We use the comprehensive study of Kim *et al.*[15] for constraints resulting from neutrino-induced and electron-induced interactions. We then add the updated constraints from e^+e^- annihilation,[16] and atomic-physics parity-nonconservation measurements.[17] In the most general case, x_L and x_R are constrained only by $x_L + x_R < 1$. However, if we impose the additional *theoretical* constraint that $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ be derivable from simple grand unification such as SO(10),[18] then we must have at least $x_L \leq x_R$, because $SU(2)_R$ must be broken at an energy scale greater than or equal to the corresponding one for $SU(2)_L$. In addition, if $U(1)_R$ breaks off from the grand-unification symmetry at a lower energy scale than $U(1)_{B-L}$, then $x_R < \frac{3}{5}(1-x_L)$. Hence we consider below both the general case $x_R \leq 1 - x_L$ as well as the more restricted case $x_L \leq x_R \leq 0.65(1-x_L)$. [We allow x_R to exceed $\frac{3}{5}(1-x_L)$ by a small amount to accomodate the unusual situation[11] where $U(1)_R$ breaks off at a higher energy than $U(1)_{B-L}$.] From Eqs. (8)–(10), we see that both ρ_1 and η can be taken to be positive whereas ρ_2 can then be of either sign. Comparing Eq. (7) with the standard-model limit where only J_{ZL} is involved, we also see that $\rho_1^2 + \rho_2^2$ must turn out to be near unity, and η less than unity.

To display our results fully, we would need a five-dimensional plot in parameter space, which we are unable to do here. Let us just summarize the main features of our analysis with regard to x_L, x_R, ρ_1, ρ_2, η and then show in Figs. 1 and 2 how M_{Z_1}, M_{Z_2}, and M_{W_1} are constrained as a result. The numbers we quote are one-standard-deviation limits for five independent variables. The parameters ρ_1 and ρ_2 satisfy $0 < \rho_1 < 1.1$ and $-1.1 < \rho_2 < 1.1$, but they are highly correlated such that $0.85 < \rho_1^2 + \rho_2^2 < 1.2$. We also find $0 < \eta < 0.5$, $0.17 < x_L < 0.28$, and $0 < x_R < 0.82$, where the only theoretical constraint used is $x_R \leq 1 - x_L$. If we take the more restricted case $x_L \leq x_R \leq 0.65(1-x_L)$, the allowed ranges for ρ_1, ρ_2, and η remain unaffected, but those for x_L and x_R are reduced to $0.205 < x_L < 0.265$ and $0.2 < x_R < 0.52$.

For given x_L, x_R, ρ_1, ρ_2, and η, we can invert Eqs. (8)–(10) to obtain A, B, and C, which in turn give us the two mass eigenvalues M_{Z_1} and M_{Z_2}. In Fig. 1, we show the region in M_{Z_2}/M_Z vs M_{Z_1}/M_Z allowed by the experimental constraints on the five phenomenological parameters.

We normalize M_{Z_1} and M_{Z_2} against the value $M_Z = 93.8$ GeV expected for the single neutral vector boson in the standard model. We find that M_{Z_2} must be greater than 200 GeV, and that M_{Z_1} must be between 83 and 116 GeV if only $x_R \leq 1 - x_L$ is assumed; M_{Z_1} is between 86 and 102 GeV if $x_L \leq x_R \leq 0.65(1 - x_L)$. In Fig. 2, we show the allowed ranges for M_{Z_1} as a function of x_L. We see the expected correlation between low x_L and high M_{Z_1}. We also show the plot $M_{W_1} = M_W(x_W/x_L)^{1/2}$, where $x_W = \sin^2\theta_W$ for the standard model, for the case of negligible mixing between W_L and W_R. With $M_W = 83.0$ GeV, we find 75 GeV$< M_{W_1} <$ 97 GeV in the most general case and 78 GeV$< M_{W_1} <$ 88 GeV for $x_L \leq x_R \leq 0.65(1 - x_L)$. (The lower bound on M_{W_2} will depend on the model. In the case of manifest left-right symmetry, it has been found that[4] $M_{W_2}/M_{W_1} > 2.8$.)

If we use the two-standard-deviation limits in our analysis, the allowed region in parameter space expands, but only by several percent. Also, if we leave out the constraint due to parity-nonconservation effects in heavy atoms, no significant additions in the allowed region are apparent. There are obviously many more interesting results to be extracted from our analysis, such as the identification of regions in parameter space with specific Higgs representations, etc. We leave this and all other discussions to another paper where we will also present details of our phenomenological analysis.

In conclusion, we have used existing experimental data to obtain highly restrictive bounds on M_{Z_1} and M_{Z_2} for the most general $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge model. Our analysis helps to define the allowed range of mass values for the possible observation of the first weak bosons. Since left-right models are on equally firm ground theoretically as the standard model, it is important to know what masses are already excluded by present data without regard to the details of the model. The results in Figs. 1 and 2 can be used for crucial tests of $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge models when the $p\bar{p}$ collider at CERN becomes capable of producing $W^{\pm}$ and Z^0 bosons at an observable rate.

This work was supported in part by the U. S. Department of Energy under Contracts No. DE-AM03-76SF00235 and No. DE-AC02-76ER00881.

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