Some phenomenological aspects of the $SU(2)_q \times SU(2)_l \times U(1)_Y$ model

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We examine the implications of the recently proposed un-unified $SU(2)_q \times SU(2)_l \times U(1)_Y$ model for the extra-gauge-boson decay widths and hadronic production rates, the $e^+e^- \to \mu^+\mu^-$ and $b\bar{b}$ asymmetries, the B-meson semileptonic branching fraction, B^0-$\bar{B}^0$ mixing, and the CP-violating parameter ϵ. A perturbative requirement gives the bound $\sin\phi \gtrsim 0.22$ on the new mixing angle. Current searches at the Fermilab Tevatron collider for an extra Z boson can place a mass bound $M_{Z_H} \gtrsim 250$ GeV. The hadronic decays of W_H, Z_H produced in $\bar{p}p$ collisions may be observable as enhancements in the two-jet invariant-mass distribution. The enhanced low-energy nonleptonic interaction improves the agreement with experiment of the predicted B-meson semileptonic branching fraction.

I. INTRODUCTION

A number of electroweak gauge theory alternatives to the standard model with expanded gauge symmetry groups containing extra U(1) and/or SU(2) factors at low energy have been extensively considered,[1] such as the left-right-symmetric model $SU(2)_L \times SU(2)_R \times U(1)$ and the E_6 superstring-inspired model $SU(2)_L \times U(1)_Y \times U(1)_\eta$. The contributions associated with the extra Z and/or W bosons modify the standard-model phenomenology at low energy and constraints can thereby be placed on the masses and couplings of the extra weak bosons. These alternative models thus provide useful yardsticks by which the success of the standard model can be measured. Recently Georgi, Jenkins, and Simmons[2] introduced an ingenious alternative electroweak gauge group $SU(2)_q \times SU(2)_l \times U(1)_Y$, in which left-handed quarks and leptons transform under separate SU(2) symmetries and right-handed fermions transform as singlets under both. This un-unified model closely duplicates the low-energy phenomenology of the standard model, even if the extra W and Z bosons are relatively light. In this paper we obtain further constraints on the parameters of the model from perturbative requirements on the couplings, the B-meson semileptonic branching fraction, and B^0-$\bar{B}^0$ mixing. We also discuss prospects for detection of the new weak bosons in $p\bar{p}$ collider experiments at CERN, Fermilab Tevatron, and Superconducting Super Collider (SSC) energies. Before proceeding we briefly review the salient aspects of the model in this section.

The additional parameters of the un-unified model are a new mixing angle ϕ and a ratio $x=(u/v)^2$ of squares of scalar vacuum expectation values (VEV's). The electroweak couplings are given in terms of ϕ as

$$g_\phi = \frac{e}{s_\phi s_W}, \quad g_l = \frac{e}{c_\phi s_W}, \quad g' = \frac{e}{C_W}, \tag{1}$$

where $s_\phi \equiv \sin\phi$, $c_\phi \equiv \cos\phi$, $s_W = \sin\theta_W$, and $c_W = \cos\theta_W$. The electroweak symmetry is spontaneously broken by the vacuum expectation values of two scalar multiplets

$$\langle \Sigma \rangle = \begin{bmatrix} u & 0 \\ 0 & u \end{bmatrix}, \quad \langle \phi \rangle = \begin{bmatrix} 0 \\ v/\sqrt{2} \end{bmatrix}. \tag{2}$$

The VEV of Σ breaks the two SU(2)'s down to the SU(2) of the standard model (SM).

The interaction Lagrangian is conveniently expressed in the bases

$$\begin{bmatrix} W_1^\pm \\ W_2^\pm \end{bmatrix} = \begin{bmatrix} c_\phi & s_\phi \\ -s_\phi & c_\phi \end{bmatrix} \begin{bmatrix} W_l^\pm \\ W_q^\pm \end{bmatrix}, \qquad Z_1 = c_W W_1^0 - s_W X, \quad Z_2 = W_2^0, \tag{3}$$

where W_l, W_q, X refer to the original gauge-boson eigenstates. The $W_1^\pm$ and Z_1 have standard-model couplings to fermions. The couplings of the $W_2^\pm$ and Z_2 are given by

$$\mathcal{L}_{W_2} = \frac{g}{\sqrt{2}} \left[\left[-\frac{s_\phi}{c_\phi} \right] \bar{\nu}\gamma_\mu l_L + \left[\frac{c_\phi}{s_\phi} \right] \bar{u}\gamma_\mu d_L \right] W_2^+ ,$$

$$\mathcal{L}_{Z_2} = \frac{g}{2} \left[\left[-\frac{s_\phi}{c_\phi} \right] (\bar{\nu}\gamma_\mu \nu_L - \bar{l}\gamma_\mu l_L) + \left[\frac{c_\phi}{s_\phi} \right] (\bar{u}\gamma_\mu u_L - \bar{d}\gamma_\mu d_L) \right] Z_2 , \tag{4}$$

where $\psi_L = \frac{1}{2}(1-\gamma_5)\psi$ denotes a left-handed spinor. The mass-squared matrices for the charged W's and the Z's are diagonalized in the bases $W_L^\pm, W_H^\pm$ and $Z_L^\pm, Z_H^\pm$ defined by the rotations

$$\begin{bmatrix} W_1^\pm \\ W_2^\pm \end{bmatrix} = \begin{bmatrix} c_\chi & s_\chi \\ -s_\chi & c_\chi \end{bmatrix} \begin{bmatrix} W_L^\pm \\ W_H^\pm \end{bmatrix} ,$$
$$\begin{bmatrix} Z_1 \\ Z_2 \end{bmatrix} = \begin{bmatrix} c_\xi & s_\xi \\ -s_\xi & c_\xi \end{bmatrix} \begin{bmatrix} Z_L \\ Z_H \end{bmatrix} , \tag{5}$$

with angles determined to be

$$\tan 2\chi = \frac{-2t_\phi}{\frac{x}{s_\phi^2 c_\phi^2} + t_\phi^2 - 1}, \quad \tan 2\xi = \frac{-2t_\phi / c_W}{\frac{x}{s_\phi^2 c_\phi^2} + t_\phi^2 - c_W^{-2}} , \tag{6}$$

where $t_\phi \equiv \tan\phi$, $c_\chi \equiv \cos\chi$, $c_\xi \equiv \cos\xi$, etc. In terms of $y = x/s_\phi^2$ the resulting masses are

$$M_{W_L, W_H} = \frac{M_W}{\sqrt{2} c_\phi} \{ y + 1 \mp [(y+1)^2 - 4yc_\phi^2]^{1/2} \}^{1/2} ,$$
$$M_{Z_L, Z_H} = \frac{M_W}{\sqrt{2} c_\phi} \{ y + 1 \mp [(y+1)^2 - 4yc_\phi^2 / c_W^2]^{1/2} \}^{1/2} , \tag{7}$$

where M_W and M_Z are the standard-model masses at the tree level. Consistency of the extra neutral-current contributions to electro-deuteron scattering, $e^+e^- \to \mu^+\mu^-$ and Bhabha-scattering measurements is achieved if s_ϕ^2 is not too large and x is not too small. The prediction for the $e^+e^- \to \mu^+\mu^-$ asymmetry at $\sqrt{s} \approx 35$ GeV, relevant to DESY PETRA measurements,[3] deviates significantly from the SM result only at $s_\phi \gtrsim 0.8$ for $M_{Z_H} = 100$ GeV ($s_\phi \gtrsim 0.95$ for $M_{Z_H} = 200$ GeV). The $e^+e^- \to \bar{b}b$ asymmetry is essentially independent of s_ϕ for $s_\phi \gtrsim 0.1$ and the deviation from the SM prediction is $\delta A_{\mathrm{FB}} \approx 0.2$ (100 GeV$/M_{Z_H})^2$. The predicted deviations from the SM results at KEK TRISTAN energies[3] are similar but somewhat smaller. Thus for $M_{Z_H} \gtrsim 200$ GeV the asymmetry measurements do not give restrictions on the model.

For the relevant condition that $x/s_\phi^2 \gg 1$ the approximate masses are

$$M_{W_L} \simeq M_W \left[1 - \frac{s_\phi^4}{2x} \right] ,$$
$$M_{Z_L} \simeq M_Z \left[1 - \frac{s_\phi^4}{2x} \right] , \tag{8}$$
$$M_{W_H} \simeq M_{Z_H} \simeq \frac{M_W \sqrt{x}}{s_\phi c_\phi} \left[1 + \frac{s_\phi^4}{2x} \right] .$$

In this large-x/s_ϕ^2 approximation the mass eigenstates W_L, W_H, Z_L, Z_H, and the states of Eq. (3) are related by

$$W_L \approx W_1 + \frac{s_\phi^3 c_\phi}{x} W_2, \quad W_H \approx W_2 - \frac{s_\phi^3 c_\phi}{x} W_L ,$$
$$Z_L \approx Z_1 + \frac{s_\phi^3 c_\phi}{x c_W} Z_2, \quad Z_H \approx Z_2 - \frac{s_\phi^3 c_\phi}{x c_W} Z_1 . \tag{9}$$

In our considerations we will approximate the mass eigenstates by the states W_1, W_2, Z_1, and Z_2.

II. TOTAL WIDTHS OF W_2, Z_2

The couplings in Eq. (4) of W_2 and Z_2 to quarks are proportional to c_ϕ/s_ϕ and can thus be potentially very large in the small-s_ϕ limit of interest. Consequently if we impose the requirement that calculations be perturbative, we can place a lower bound on s_ϕ. A reasonably conservative demand is that the total width not exceed one-half the mass, $\Gamma \lesssim \frac{1}{2}M$, since it seems unlikely that a gauge boson with a width $\Gamma \gtrsim \frac{1}{2}M$ would be experimentally identifiable as a resonance state.

The W_H and Z_H partial widths are

$$\Gamma(W_H \to l\nu) = 2\Gamma(Z_H \to \nu_l \bar{\nu}_l) = 2\Gamma(Z_H \to l\bar{l}) = \left[\frac{s_\phi}{c_\phi} \right]^2 \frac{G_F M_W^2 M_H}{6\sqrt{2}\pi} ,$$
$$\Gamma(W_H \to u\bar{d}) \simeq 2\Gamma(Z_H \to u\bar{u}) = 2\Gamma(Z_H \to d\bar{d}) = 3 \left[\frac{c_\phi}{s_\phi} \right]^2 \frac{G_F M_W^2 M_H}{6\sqrt{2}\pi} , \tag{10}$$

where quark mixing has been neglected. In the approximation that decays to t quarks are included without phase-space suppression, the total widths are

$$\Gamma(W_H \to \text{all}) = \Gamma(Z_H \to \text{all}) = \left[\left[\frac{s_\phi}{c_\phi} \right]^2 + 3 \left[\frac{c_\phi}{s_\phi} \right]^2 \right] \frac{G_F M_W^2 M_{W_H}}{2\sqrt{2}\pi} . \tag{11}$$

Figure 1 illustrates $\Gamma(Z_H \to \text{all})/M_{Z_H}$ versus s_ϕ for $M_{Z_H} = 200$ GeV and $m_t = 80$ GeV. The total width for W_H is essentially the same. For small s_ϕ, the relation be-

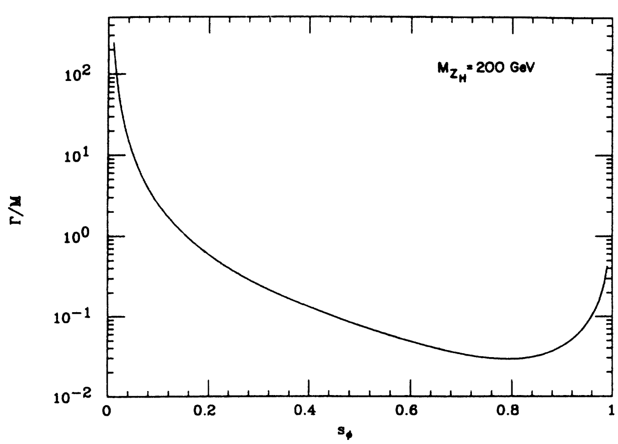


FIG. 1. Total width for Z_H decays (divided by M_{Z_H}) versus s_ϕ for $M_{Z_H} = 200$ GeV and $m_t = 80$ GeV. The total width for W_H is essentially the same.

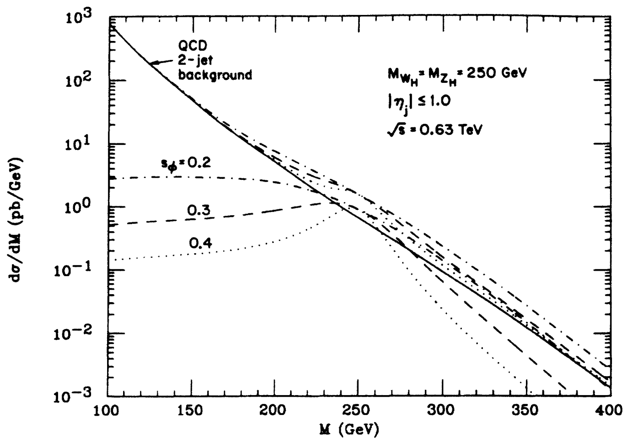


FIG. 2. Invariant two-jet mass distribution in $\bar{p}p$ collisions at $\sqrt{s}=0.63$ TeV from W_H, Z_H production and the QCD background, taking $M_{W_H}=M_{Z_H}=250$ GeV. The mixing values $s_\phi=0.2$, 0.3, and 0.4 are illustrated.

In Figs. 2 and 3 the independent parameters of the model are taken to be s_ϕ and M_{Z_H}. Then for $x/s_\phi^2 \gg 1$ the value of x is determined by

$$x \simeq s_\phi^2 (M_{Z_H}/M_W)^2 .$$

The nonleptonic interactions in the model are enhanced by a factor of $1+1/x$ and consequently values of $x \lesssim 1$ are probably excluded. This in turn places a constraint

$$s_\phi \gtrsim M_W/M_{Z_H} .$$

Values of $s\phi$ that do not satisfy this condition are therefore unlikely to be of physical interest.

At the Tevatron the presence of W_H and Z_H with mass $M=300$ GeV and mixing $s_\phi=0.4$ would give a cross-section enhancement over a range $\Delta M=30$ GeV of $\Delta\sigma=124$ pb compared to a QCD background cross section of $\sigma_{\rm QCD}=600$ pb in this mass range. With the present 4.7-pb^{-1} of luminosity this effect corresponds to $\Delta N=580$ events over the QCD expectation of $N=2820$ events. Thus the model parameters can soon be con-

tween the total width of W_H and the standard W width becomes

$$\Gamma(W_H \to \text{all})/M_{W_H} \approx \left[\frac{1}{s_\phi}\right]^2 \Gamma(W \to \text{all})/M_W . \qquad (12)$$

Our perturbative criteria that

$$\frac{\Gamma(W_H \to \text{all})}{M_{W_H}} \lesssim \frac{1}{2}$$

gives the s_ϕ bound

$$s_\phi \gtrsim 0.22 \qquad (13)$$

for which the effective quark coupling strength is $(c_\phi/s_\phi)g \sim 5g$. By the same criteria we also must deduce that $s_\phi \lesssim 0.99$.

III. SIGNALS OF W_H, Z_H AT HADRON COLLIDERS

The W_H and Z_H bosons can be produced at $\bar{p}p$ and pp colliders. The dominant hadronic decays of these broad resonances can give enhancements over the two-jet QCD background. The UA2 Collaboration[4] has extracted the $W,Z \to q\bar{q}$ signal at the CERN $\bar{p}p$ collider, and it may also be possible to detect the hadronic decay signal of W_H and Z_H. Figures 2 and 3 show representative effects of W_H, Z_H production on the two-jet invariant-mass distribution at CERN and Tevatron $\bar{p}p$ collider energies, respectively; Fig. 2 illustrates the cases $M_{W_H}=M_{Z_H}=250$ GeV for $s_\phi=0.2$, 0.3, and 0.4 at $\sqrt{s}=0.63$ TeV. Figure 3 illustrates the cases $M_{W_H}=M_{Z_H}=300$, 400 GeV for the same s_ϕ values at $\sqrt{s}=1.8$ TeV. In these two illustrations the interference between the QCD amplitude and the real part of the resonant amplitude has been neglected.

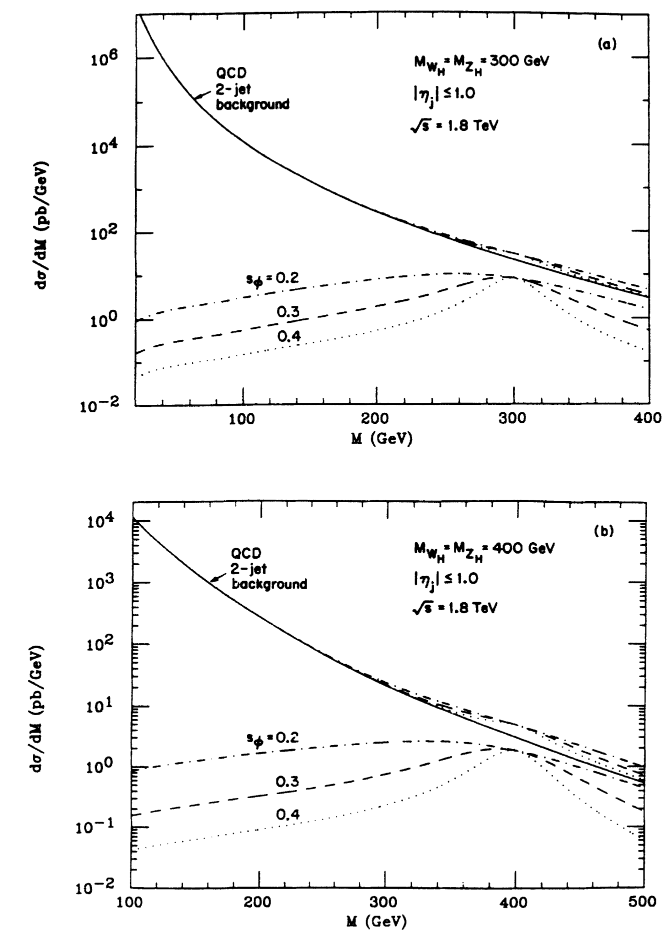


FIG. 3. Invariant two-jet mass distribution in $\bar{p}p$ collisions at $\sqrt{s}=1.8$ TeV illustrating the effects of W_H, Z_H production with mass (a) 300 GeV and (b) 400 GeV for $s_\phi=0.2$, 0.3, and 0.4.

strained further by data from the Collider Detector Facility (CDF) experiment[5] at Fermilab in which the m_{jj} distribution is measured accurately out to $m_{jj} \approx 600$ GeV.

At SSC energies the QCD amplitude is orders of magnitude larger than the W_H, Z_H hadronic signals due to the high gluon luminosity. Thus the Tevatron energy is much better suited to a search for W_H, Z_H in the hadronic decay modes.

The leptonic signals of W_H, Z_H resonances produced in hadronic collisions are promising, even though the leptonic branching fractions

$$B(W_H \to l\nu) = 2B(Z_H \to l\nu) = \frac{t_\phi^4}{9+3t_\phi^2} \tag{14}$$

are small for small s_ϕ, due to the larger production cross section resulting from enhanced W_H, Z_H couplings to quarks at small s_ϕ. Figures 4, 5, and 6 present expectations at the CERN, Tevatron, and SSC colliders, respectively, for (a) $\sigma B(W_H \to e\nu)$ and (b) $\sigma B(Z_H \to ee)$ signals versus s_ϕ for various W_H and Z_H mass choices. The preliminary CDF limit[6] $\sigma B < 0.85$ pb on an extra Z boson gives a bound $M_{Z_H} \gtrsim 250$ GeV for the s_ϕ bound in Eq. (13).

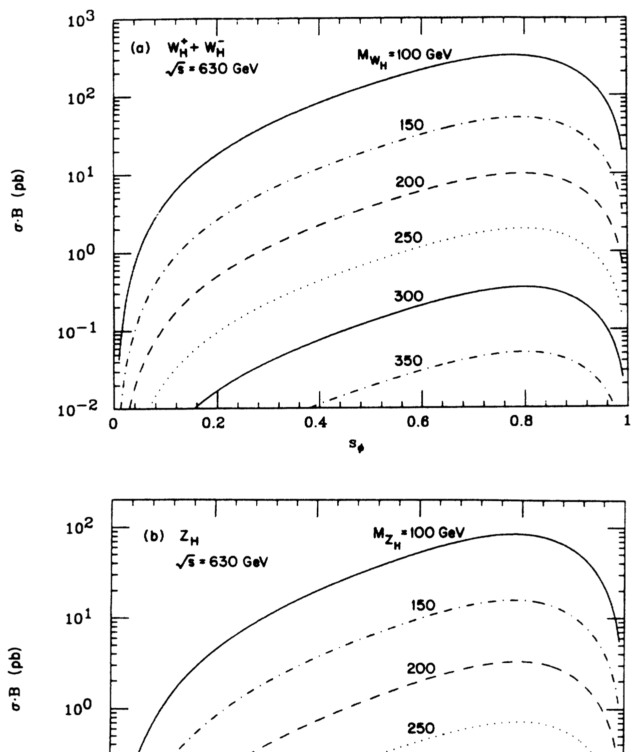


FIG. 4. Cross sections for (a) $W_H \to e\nu$, (b) $Z_H \to e\bar{e}$ production for the CERN $p\bar{p}$ collider energy of $\sqrt{s} = 0.63$ TeV.

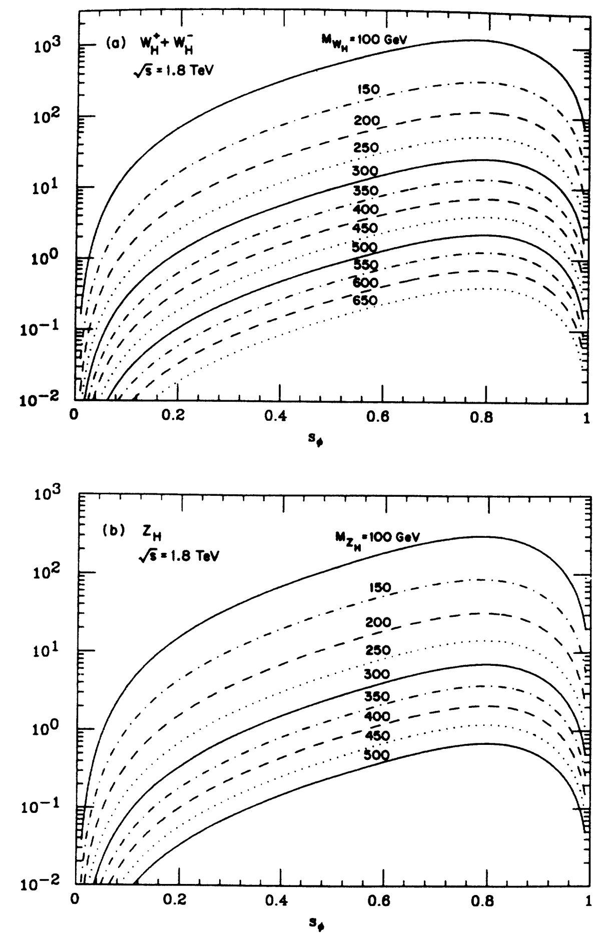


FIG. 5. Cross sections for (a) $W_H \to e\nu$, (b) $Z_H \to e\bar{e}$ production for the Tevatron $p\bar{p}$ collider energy of $\sqrt{s} = 1.8$ TeV.

The $M_T(e\nu)$ transverse mass and $M(ee)$ invariant-mass distributions which would result at the Tevatron energy for $M_{W_H} = M_{Z_H} = 300$ GeV are shown in Fig. 7.

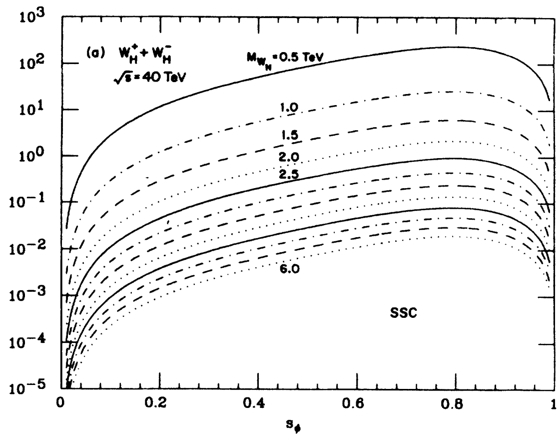
IV. B-SEMILEPTONIC BRANCHING FRACTION

The effective charged-current Lagrangian at low energy for this model is of the form

$$-\mathcal{L}_{\text{eff}}^{\text{CC}} = \frac{2}{v^2}\left[(j_q + j_l)^2 + \frac{1}{x} j_q^2\right], \tag{15}$$

where the $(j_q + j_l)^2$ term duplicates the standard-model interactions. The nonleptonic weak interactions are enhanced relative to the standard model by a factor of $1 + 1/x$. A direct consequence is a smaller predicted value for the semileptonic branching fraction for B mesons.

In the spectator decay model, the $b \to c$ transitions are known to be dominant over $b \to u$ transitions. Estimating[7] the phase-space suppressions to be $\Gamma(b \to c\tau\nu)/\Gamma(b \to ce\nu) \simeq 0.2$ and $\Gamma(b \to c\bar{c}s)/\Gamma(b \to c\bar{u}d) \simeq 0.14$, the B-meson semileptonic branching fraction is



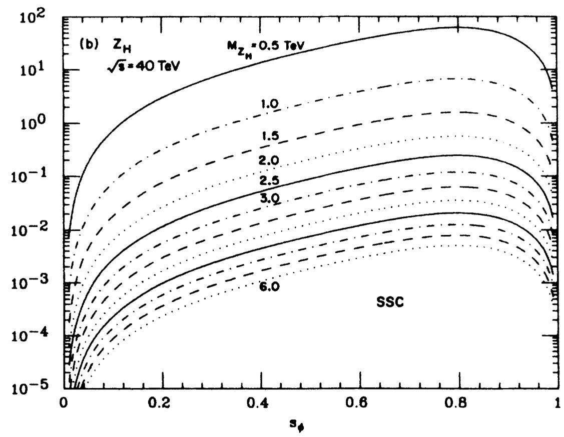


FIG. 6. Cross sections for (a) $W_H \to e\nu$, (b) $Z_H \to e\bar{e}$ production for the SSC pp collider energy of $\sqrt{s}$ = 40 TeV.

$$B(B \to cl\nu) = \frac{1}{2.2 + 1.14(2C_+^2 + C_-^2)\left[1+\frac{1}{x}\right]^2} , \qquad (16)$$

where

$$\tfrac{1}{3}(2C_+^2 + C_-^2) \simeq \tfrac{1}{3}[2(1.38)^2 + (0.85)^2] \simeq 1.51$$

is the QCD enhancement of the nonleptonic decay rate.[8] Equation (16) with $1/x = 0$ yields the standard-model prediction $B = 13.6\%$ with an estimated[9] theoretical error of $\pm 1.6\%$. Recent measurements[10] from the ARGUS and CLEO Collaborations give $B = (10.0 \pm 0.8)\%$ and $B = (9.3 \pm 2.4 \pm 0.9)\%$, respectively. For $x = 5$ the predicted value is $B = 10.4\%$. Thus the additional nonleptonic interaction in Eq. (15) can resolve the discrepancy with experiment with an appropriate value of x.

V. B^0-$\bar{B}^0$ MIXING

The standard-model calculation of B^0-$\bar{B}^0$ mixing[11] is modified by the presence of additional box diagrams involving W_H. Only the contributions of box diagrams with t-quark exchanges are considered here, since this is an excellent approximation when the top quark is heavy [the preliminary CDF experimental limit[12] is $m_t \geq 78$ GeV assuming that $B(t - bl\nu) \simeq 0.1$]. The mixing amplitude resulting from the diagram with two W_L exchanges is given by

$$A_{W_L W_L} = G(x) \equiv x[\tfrac{1}{4} + \tfrac{9}{4}(1-x)^{-1} - \tfrac{3}{2}(1-x)^{-2}] - \frac{3}{2}\left[\frac{x}{1-x}\right]^3 \ln x , \qquad (17)$$

where $x = m_t^2/M_{W_L}^2$; an overall constant factor is omitted since we are only interested in relative contributions. Similarly the diagram with two W_H exchanges has the amplitude

$$A_{W_H W_H} = \frac{1}{t_\phi^4} G(y) , \qquad (18)$$

where $y = m_t^2/M_{W_H}^2$. The net amplitude from the $W_L W_H$ exchange diagrams (including unphysical-Higgs-boson exchanges) in terms of $r = M_{W_H}^2/M_{W_L}^2$ are found to be

$$A_{W_L W_H} = -\frac{2}{t_\phi^2}[(I_1 - 2I_3 + 1) + 2(x^2/r)I_2 - \tfrac{1}{4}(x^2/r)I_1] , \qquad (19)$$

where the integrals I_1, I_2, and I_3 have the expressions

$$I_1 \equiv \int_0^\infty \frac{z^2 dz}{(z-x)^2(z-1)(z-r)} = \frac{x}{(x-r)(1-x)} + \frac{r^2 \ln r}{(1-r)(x-r)^2} + \frac{x^2 + x^2 r - 2xr}{(x-r)^2(1-x)^2} \ln x , \qquad (20a)$$

$$I_2 \equiv \int_0^\infty \frac{z\, dz}{(z-x)^2(z-1)(z-r)} = \frac{1}{(x-r)(1-x)} + \frac{r \ln r}{(x-r)^2(1-r)} + \frac{x^2 - r}{(x-r)^2(1-x)^2} \ln x , \qquad (20b)$$

$$I_3 \equiv \int_0^\infty \frac{z\, dz}{(z-x)(z-1)(z-r)} = \left[1 + \frac{x(1+r-2x)}{2(1-x)(r-x)}\right]\frac{\ln r}{r} - \frac{x(2\ln x - \ln r)}{2(1-x)(r-x)} . \qquad (20c)$$

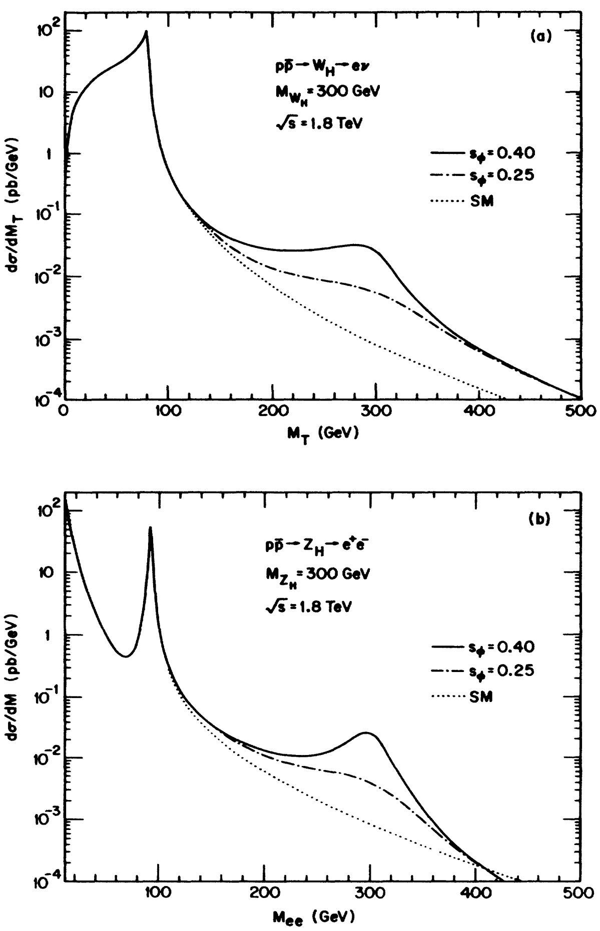


FIG. 7. Predicted transverse mass $M_T(e\nu)$ and invariant-mass $M(e\bar{e})$ distributions at $\sqrt{s}=1.8$ TeV for $M_{W_H}=M_{Z_H}=300$ GeV.

Using these expressions we evaluate the ratio

$$R=\frac{A_{W_L W_L}+A_{W_H W_H}+A_{W_L W_H}}{A_{W_L W_L}} \tag{21}$$

which is the enhancement of $\delta m/\Gamma$ for B^0-$\bar{B}^0$ mixing relative to the standard-model expectation. Figure 3 shows R versus s_ϕ with $M_H=100$, 200, 300, 400, and 1000 GeV for (a) $m_t=90$ GeV and (b) $m_t=120$ GeV. Since large enhancements ($R \gtrsim 5$) are very common at small s_ϕ, large regions of the s_ϕ, M_{W_H} parameter space can be excluded on this basis; however there are particular M_{W_H}, s_ϕ parameter choices such as $M_{W_H} \simeq 300$ GeV, $s_\phi \sim 0.2$ for which $R \sim 1$.

The predicted enhancement for the CP-violating parameter ϵ relative to the standard model will be the same as for $\delta m/\Gamma$ is the contributions to ϵ are dominated by the top-quark-exchange box diagrams.

VI. THEORETICAL CONSTRAINTS

As it stands, the $SU(2)_q \times SU(2)_l \times U(1)_Y$ model is incomplete in at least two ways. First, the minimal SM fermion content leads to anomalies which must be canceled by the addition of new fermions. Second, in order to generate fermion masses the scalar sector of the model must be extended.

The fermions that are added to cancel the anomalies must be such that the usual SM Z does not couple in an off-diagonal manner leading to flavor-changing neutral currents (FCNC's). If we demand that Z_H also couples in a way which is flavor diagonal, then anomaly cancellation becomes highly nontrivial. Here we will consider the simpler case where Z_H exchange leads to FCNC which are suppressed by the large value of M_H, while all SM Z couplings are flavor diagonal. (Note that the addition of mirror fermions, while canceling anomalies, leads to significant FCNC for *both* the Z and Z_H.) To ensure that absences of FCNC couplings to the SM Z we need only assume that the new fermions that we add to cancel the anomalies have the same values of $T_{3l}+T_{3q}$ as do the SM fermions of the same charge. (T_{3l} and T_{3q} may be *separately* different from those of the SM fermions.) The nontrivial anomaly cancellation requirements are

$$\mathrm{Tr}\left[\frac{Y}{2}\right]^3=\mathrm{Tr}\left[T_{3q}^2\frac{Y}{2}\right]=\mathrm{Tr}\left[T_{3l}^2\frac{Y}{2}\right]=0\;, \tag{22}$$

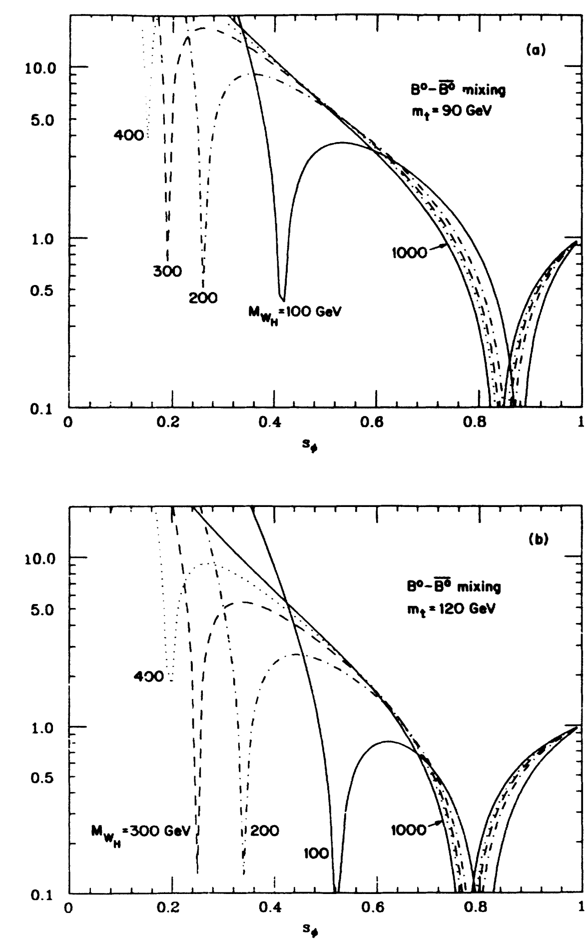


FIG. 8. Ratio R of the B^0-$\bar{B}^0$ mixing amplitude $\delta m/\Gamma$ to the standard-model result.

where we have assumed that *all* fermions transform as either singlets or doublets under $SU(2)_{l,q}$. Now consider the new color-triplet fermions

$$\begin{pmatrix} x_1 \\ x_2 \end{pmatrix}_L^a \text{ and } (x_{1,2}^a)_R, \quad Q(x_{1,2}) = \tfrac{2}{3}, -\tfrac{1}{3} \tag{23}$$

transforming as a doublet and singlets under $SU(2)_l$ and the new color-singlet fermions

$$\begin{pmatrix} y_1 \\ y_2 \end{pmatrix}_L \text{ and } (y_{1,2})_R, \quad Q(y_{1,2}) = 0, -1 \tag{24}$$

transforming as a doublet and singlets under $SU(2)_q$ ($a = 1,2,3$ is a color index). This means, for example, that x_{1L} transforms as a member of a doublet under $SU(2)_l$ but is an $SU(2)_q$ singlet while y_{1L} transforms as a member of a doublet under $SU(2)_q$ but is an $SU(2)_l$ singlet. All right-handed fields are singlets under both groups. Note that while the sum $T_{3l} + T_{3q}$ for x_{1L} is the same as that for u_L, T_{3l} and T_{3q} are separately distinct. These new fermions then lead to anomaly cancellation as well as flavor-conserving couplings for the SM Z. However, Z_H has FCNC interactions when the new and ordinary SM fermions mix. Once intergenerational mixing is turned on, the fact that the new and SM fermions have different values of T_{3l} and T_{3q} will lead to generic flavor-changing Z_H couplings of the form, e.g.,

$$\mathcal{L} \sim \frac{1}{4} g \left(\frac{s_\phi}{c_\phi} \right) \theta^2 [\bar{e} \gamma_\mu (1 - \gamma_5) \mu + \text{H.c.}] Z_H , \tag{25}$$

where θ describes a typical mixing between the new fermions and the ones already contained in the SM. From this we estimate that

$$R = \frac{\Gamma(\mu \to 3e)}{\Gamma(\mu \to e \nu \bar{\nu})} \approx \left(\frac{s_\phi}{c_\phi} \right)^4 \theta^4 \left(\frac{M_{W_L}}{M_{Z_H}} \right)^4 . \tag{26}$$

For R to be less than the experimental bound[13] of 10^{-12} requires that

$$\left(\frac{s_\phi}{c_\phi} \right) \theta \frac{M_{W_L}}{M_{Z_H}} \lesssim 10^{-3} . \tag{27}$$

For $s_\phi \simeq 0.3$ and $\theta \approx 0.05$ this would imply that $M_{Z_H} \gtrsim 1.4$ TeV. It has been argued that a similar M_H bound occurs when the addition of new scalars to accommodate fermion mass generation is considered.[14]

The theoretical constraints discussed in this section are based on particular extensions to make the model complete. It is possible that there exist different extensions that allow lower W_H, Z_H masses. For that reason direct experimental searches at the Tevatron, as discussed in Sec. III, should still be pursued.

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