Phenomenology of the B-$3L_\tau$ gauge boson

Ernest Ma
Department of Physics, University of California, Riverside, California 92521

D. P. Roy
Tata Institute of Fundamental Research, Mumbai 400 005, India

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Assuming the existence of a gauge boson X which couples to B-$3L_\tau$, we discuss the present experimental constraints on g_X and m_X from $Z\to l^+l^-$ and $Z\to\bar{f}fX(f=q,\nu_\tau,\tau)$. We also discuss the discovery potential of X at hadron colliders through its decay into $\tau^+\tau^-$ pairs. In the scenario where all three charged leptons (and their neutrinos) mix, lepton flavor nonconservation through X becomes possible and provides another experimental probe into this hypothesis. [S0556-2821(98)07319-6]

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I. INTRODUCTION

In the minimal standard model of quarks and leptons, neutrinos (ν_e, ν_μ, ν_τ) appear only as members of left-handed doublets and there is a single Higgs scalar doublet. Hence neutrinos are massless and each lepton number (L_e, L_μ, L_τ) is separately conserved at the classical level as is baryon number (B). However, only the linear combination B-L_e-L_μ-L_τ remains conserved at the quantum level [1], whereas the corresponding U(1) is still anomalous and cannot be gauged. Actually, one of the three lepton number differences (L_e-L_μ, L_e-L_τ, L_μ-L_τ) is anomaly-free and could be gauged [2]. On the other hand, this particular extension of the standard model would not shed much light on the question of neutrino mass. After all, there is now a host of experimental evidence for neutrino oscillations and that can be explained most naturally if neutrinos have masses and mix with one another. The canonical way of doing this is to add three right-handed neutrino singlets with large Majorana masses so that ν_e, ν_μ, and ν_τ all obtain small seesaw masses. Such an extension of the standard model allows B-L to be gauged [3]. (Since the three lepton families now mix, it makes sense to consider only one lepton number, i.e., $L=L_e+L_\mu+L_\tau$.)

Suppose we add only *one* right-handed neutrino singlet and pair it with ν_τ. Then the symmetry $B-3L_\tau$ can be gauged [4]. Just as B-L may originate [5] from $\mathrm{SU}(4)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R$, the breaking of $\mathrm{SU}(10)\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_{Y'}$ to the standard gauge group by way of SU(9) leads naturally [4] to B-$3L_\tau$. This recent discovery opens up a possible rich phenomenology associated with the B-$3L_\tau$ gauge boson which we call X. The key observation is that X is not constrained by present experimental data to be very heavy because it does not couple to leptons of the first and second families.

In Sec. II we describe the B-$3L_\tau$ model and show how all anomalies are canceled. In Sec. III we determine the present experimental constraints on the mass and coupling of X. They come chiefly from the nonobservation of X in the decay of Z and the agreement with e-μ-τ universality in Z decays. In Sec. IV we consider the production of X at hadron colliders and the prospect for its detection at the Tevatron and at the LHC (Large Hadron Collider). In Sec. V we study how ν_e and ν_μ may acquire radiative masses and mix with ν_τ which has a tree-level seesaw mass. We present two scenarios, one with lepton-flavor-changing couplings for X and one without. In Sec. VI we discuss the possible manifestations of lepton flavor nonconservation in the first scenario. We also explain how the one-loop quark flavor nonconservation is naturally suppressed in this model. Finally in Sec. VII we have some concluding remarks.

II. STRUCTURE OF THE B-$3L_\tau$ GAUGE MODEL

Consider the gauge group $\mathrm{SU}(3)_C\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y\times\mathrm{U}(1)_X$, where the extra U(1) refers to the $B-3L_\tau$ symmetry. The quarks and leptons transform thus as follows:

$$\begin{pmatrix} u_i \\ d_i \end{pmatrix}_L \sim (3,2,1/6;1/3), \quad u_{iR}\sim(3,1,2/3;1/3), \quad d_{iR}\sim(3,1,-1/3;1/3); \tag{1}$$

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L \sim (1,2,-1/2;0), \quad e_R,\mu_R\sim(1,1,-1;0); \tag{2}$$

$$\begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L \sim (1,2,-1/2;-3), \quad \tau_R\sim(1,1,-1;-3), \quad \nu_{\tau R}\sim(1,1,0;-3). \tag{3}$$

In the above, only one right-handed neutrino singlet, i.e., $\nu_{\tau R}$, has been added to the minimal standard model. Since the number of $\mathrm{SU}(2)_L$ doublets remains even (it is in fact unchanged), the global SU(2) chiral gauge anomaly [6] is absent. Since the quarks and leptons are chosen to transform vectorially under the new $\mathrm{U}(1)_X$, the mixed gravitational-gauge anomaly [7] is also absent. The various axial-vector anomalies [8] are canceled as well. The $[\mathrm{SU}(3)]^2\mathrm{U}(1)_X$ and

$[\mathrm{U}(1)_X]^3$ anomalies are automatically zero because of the vectorial nature of SU(3) and $\mathrm{U}(1)_X$. The remaining conditons are satisfied as follows:

$$[\mathrm{SU}(2)]^2\mathrm{U}(1)_X:\ (3)(3)(1/3)+(-3)=0; \tag{4}$$

$$[\mathrm{U}(1)_X]^2\mathrm{U}(1)_Y:\ (3)(3)(1/3)^2[2(1/6)-(2/3)-(-1/3)]+(-3)^2[2(-1/2)-(-1)]=0; \tag{5}$$

$$[\mathrm{U}(1)_Y]^2\mathrm{U}(1)_X:\ (3)(3)[2(1/6)^2-(2/3)^2-(-1/3)^2](1/3)+[2(-1/2)^2-(-1)^2](-3)=0. \tag{6}$$

The minimal scalar content of this model consists of just the usual doublet

$$\begin{pmatrix}\phi^+\\ \phi^0\end{pmatrix}\sim(1,2,1/2;0) \tag{7}$$

and a neutral singlet

$$\chi^0\sim(1,1,0;6) \tag{8}$$

which couples to $\nu_{\tau R}\nu_{\tau R}$. As the former acquires a nonzero vacuum expectation value, the electroweak gauge symmetry $\mathrm{SU}(2)_L\times\mathrm{U}(1)_L$ breaks down to $\mathrm{U}(1)_{em}$, whereas $\langle\chi^0\rangle\neq0$ breaks $\mathrm{U}(1)_X$. The resulting theory allows $\nu_{\tau L}$ to obtain a seesaw mass and retains B as an additively conserved quantum number and L_τ as a multiplicatively conserved quantum number. The two other neutrinos, i.e., ν_e and ν_μ, are massless in this minimal scenario and cannot mix with ν_τ. In fact, L_e and L_μ are still separately conserved, and L_e-L_μ can still be gauged at this point. To obtain a phenomenologically interesting neutrino sector, i.e., to accommodate present neutrino-oscillation data, we will consider later two scenarios for extending the scalar sector to allow ν_e and ν_μ to acquire radiative masses and to mix with ν_τ.

III. CONSTRAINTS ON X FROM Z DECAY

Since the B-$3L_\tau$ gauge boson X does not couple to e or μ or their corresponding neutrinos, there is no direct phenomenological constraint from the best known high-energy physics experiments, such as e^+e^- annihilation, deep-inelastic scattering of e or μ or ν_μ on nuclei, or the observation of e^+e^- or $\mu^+\mu^-$ pairs in hadronic collisions. Although X does contribute to purely hadronic interactions, its presence is effectively masked by the enormous background due to quantum chromodynamics (QCD). However, unlike the case of a gauge boson coupled only to baryon number [9], X also couples to L_τ. Hence X may decay into $\tau^+\tau^-$ or $\overline{\nu}_\tau\nu_\tau$ and be detected that way if it is produced.

The mass and coupling of X are constrained by present experimental data (from the CERN, e^+e^- collider LEP, mainly) in two important ways. The first is direct production through Z decay:

$$Z\to(\overline{q}q,\tau^+\tau^-,\overline{\nu}_\tau\nu_\tau)+X,\quad \text{then } X\to(\overline{q}q,\tau^+\tau^-,\overline{\nu}_\tau\nu_\tau). \tag{9}$$

This applies of course only to $m_X<M_Z$. The second is through its radiative contribution to $Z\to(\tau^+\tau^-,\overline{\nu}_\tau\nu_\tau)$ which breaks e-μ-τ universality.

Because of the B-$3L_\tau$ gauge symmetry, the branching fractions of X to $\tau^+\tau^-$ and $\overline{\nu}_\tau\nu_\tau$ are very substantial. Assuming that $\nu_{\tau R}$ and the t quark are too heavy to be decay products of X, and using the parton model as a crude approximation, the branching fractions of X are

$$B(X\to\tau^+\tau^-)=54/91=0.59, \tag{10}$$

$$B(X\to\overline{\nu}_\tau\nu_\tau)=27/91=0.30, \tag{11}$$

$$B(X\to\overline{q}q)=10/91=0.11. \tag{12}$$

Consider $Z\to\overline{f}fX$, where $f=\tau,\nu_\tau,q$. In the center-of-momentum frame, let E_1 and E_2 be the energies of f and $\overline{f}$, and θ the angle between their directions. Then the square of the amplitude averaged over the polarizations of Z is easily calculated to be

$$|\overline{\mathcal{M}}|^2=(B-3L_\tau)^2g_X^2\left(\frac{g_L^2+g_R^2}{2}\right)(16E_1E_2)\times\left\{(1+\cos\theta)\left[\frac{1}{(M_Z-2E_1)^2}+\frac{1}{(M_Z-2E_2)^2}\right]+\frac{4(1-\cos\theta)}{(M_Z-2E_1)(M_Z-2E_2)}\times\left[1-\frac{E_1+E_2}{M_Z}+\frac{E_1E_2(1-\cos\theta)}{M_Z^2}\right]\right\}, \tag{13}$$

where

$$g_L=(g/\cos\theta_W)(I_{3L}-\sin^2\theta_W Q)$$

and

$$g_R=(g/\cos\theta_W)(-\sin^2\theta_W Q)$$

are the standard-model couplings of f and $\overline{f}$ to Z.

Of the 9 possible final-state combinations of Z decaying into X, two are very amenable to experimental detection: i.e.,

$$Z\to\overline{q}q+X,\quad \text{then } X\to\overline{\nu}_\tau\nu_\tau, \tag{14}$$

$$Z\to\overline{\nu}_\tau\nu_\tau+X,\quad \text{then } X\to\overline{q}q. \tag{15}$$

Both result in 2 jets plus missing energy. This channel has been widely investigated for the Higgs-boson search at LEP-I. The decay process (14) resembles the Higgs signal for its invisible decay into majorons [10,11], while Eq. (15) resembles the signal for its SM (standard model) decay. The total ALEPH data from LEP-I, corresponding to 4.5 million hadronic Z events, show no events in the 2 jets plus missing energy channel after the selection cuts [12]. We have analyzed these data in terms of the decay processes (14), (15) using a parton level Monte Carlo program. The program was earlier shown to reproduce the efficiencies of these selection cuts very well in the context of the SM Higgs signal [10];

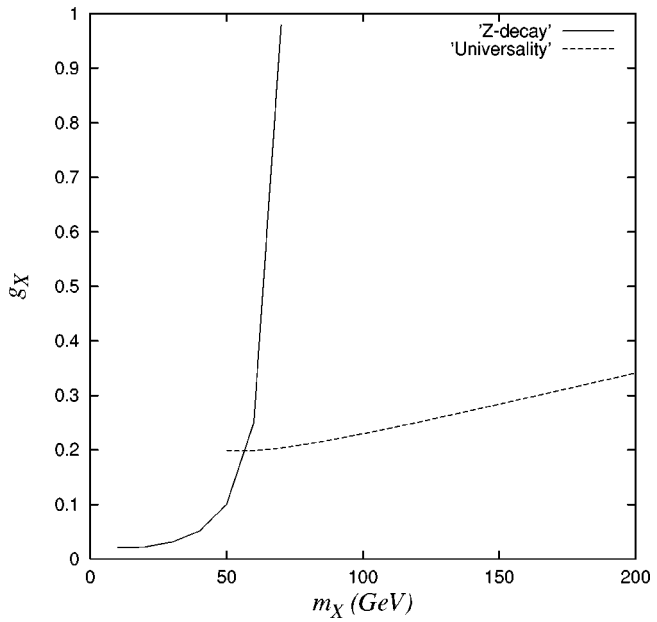

FIG. 1. The 95% C.L. upper limits on the coupling of the X boson as functions of its mass coming from Z-decay and universality constraint using LEP-I data.

and it is expected to work equally well here. In the absence of any candidate events, the 95% C.L. (confidence-level) limit on the signal corresponds to 3 events after the selection cuts. The resulting upper limit on g_X is shown as a function of m_X in Fig. 1. One gets a stringent limit on g_X (<0.1) for $m_X<50$ GeV, where the signal has a reasonable efficiency of $\sim 40\%$. But it deteriorates rapidly for $m_X \geqslant 70$ GeV, where the efficiency for the dominant process (15) goes down due to a low missing energy.

Three other final states are also important: i.e.,

$$Z\to\tau^+\tau^- + X, \quad \text{then } X\to\bar{\nu}_\tau\nu_\tau, \tag{16}$$

$$Z\to\bar{\nu}_\tau\nu_\tau + X, \quad \text{then } X\to\tau^+\tau^-, \tag{17}$$

$$Z\to\tau^+\tau^- + X, \quad \text{then } X\to\tau^+\tau^-. \tag{18}$$

The first two have two τ's plus missing energy in the final state, for which the combined decay rate is about 5 times as large as that of Eqs. (14), (15). Unfortunately the bulk of the LEP-I data in this channel have not been processed, since $H\to\tau\tau$ is not an important decay channel for the Higgs mass range of current interest. The only useful data we could find come from an early Higgs search program of ALEPH, corresponding to <0.2 million hadronic Z events [13]. The efficiency factor for this channel is slightly less than that for Eqs. (14), (15). Nonetheless one expects a factor of ~ 2 improvement in the g_X limit if the full ALEPH data in this channel are analyzed. The last process (18) corresponds to a 4τ channel and has twice as large a decay rate as Eqs. (14), (15). The detection efficiency for this channel has been estimated to be about 30% for the SM background, where the extra pair of τ's come from a virtual γ/Z [14]. The size of this background for the total data sample of 4.5 million hadronic Z events is estimated to be 2–3 events, which can be subtracted out. Assuming a similar detection efficiency for the signal process (18), one expects to get at least as good a limit on g_X from this channel as from Eqs. (14), (15).

The second constraint on g_X and m_X comes from the observed universality of $Z\to l^+l^-$ decays. Since the one-loop radiative correction of the $Z\tau^+\tau^-$ vertex has an extra contribution from X, a small deviation from e-μ-τ universality is expected. From the precision measurements [15] at the Z resonance, i.e.,

$$\Gamma_e = 83.94 \pm 0.14 \text{ MeV}, \quad \Gamma_\mu = 83.84 \pm 0.20 \text{ MeV}, \quad \Gamma_\tau = 83.68 \pm 0.24 \text{ MeV}, \tag{19}$$

and adding 0.19 MeV to Γ_τ to adjust for the kinematical correction due to m_τ, we find the deviation of Γ_τ from the average of Γ_e and Γ_μ to be bounded at 95% C.L. as follows:

$$\Delta\Gamma_\tau/\Gamma_{e,\mu} < 0.006. \tag{20}$$

Let $\delta \equiv m_X^2/M_Z^2$, then the one-loop radiative correction to $Z\to\tau^+\tau^-$ from X exchange is given by [9]

$$\frac{\Delta\Gamma_\tau}{\Gamma_\tau} = \frac{9g_X^2}{8\pi^2}F_2(\delta), \tag{21}$$

where it is well-known that

$$F_2(\delta) = -2\left\{\frac{7}{4}+\delta+\left(\delta+\frac{3}{2}\right)\ln\delta+(1+\delta)^2\left[\mathrm{Li}_2\left(\frac{\delta}{1+\delta}\right)+\frac{1}{2}\ln^2\left(\frac{\delta}{1+\delta}\right)-\frac{\pi^2}{6}\right]\right\}. \tag{22}$$

In the above, $\mathrm{Li}_2(x) = -\int_0^x (dt/t)\ln(1-t)$ is the Spence function. Using the experimental bound of Eq. (20) we show in Fig. 1 the upper limit (dashed line) on g_X as a function of m_X. Since the function F_2 decreases only slowly as δ increases, we find that the upper limit on g_X increases from 0.22 at $m_X = M_Z$ to only 0.32 at $m_X = 2M_Z$.

Assuming $g_1 \simeq 0.35$ to represent the typical size of a U(1) gauge coupling, we see that the universality limit of g_X is a fairly significant result. However it does not rule out any range of m_X. On the other hand, the Z-decay limit seems to disfavor $m_X \leqslant 40$ GeV, since the corresponding limit on g_X ($\leqslant 0.05$) is an order of magnitude smaller than g_1. For completeness we mention that the invisible width of the Z puts a constraint of $\Delta\Gamma_{\nu_\tau}/\Gamma_{\nu_\tau} < 0.015$ which is weaker than the bound of Eq. (20).

IV. PRODUCTION AND DETECTION OF X AT HADRON COLLIDERS: PRESENT CONSTRAINT AND FUTURE PROSPECT

The large branching fraction of the $X\to\tau\tau$ decay can be exploited to search for X in the $\tau\tau$ channel at hadron colliders. We shall consider the leptonic decay of one τ and hadronic decay of the other, resulting in a $l\tau$ final state. Recently the Collider Detector at Fermilab (CDF) Collaboration have

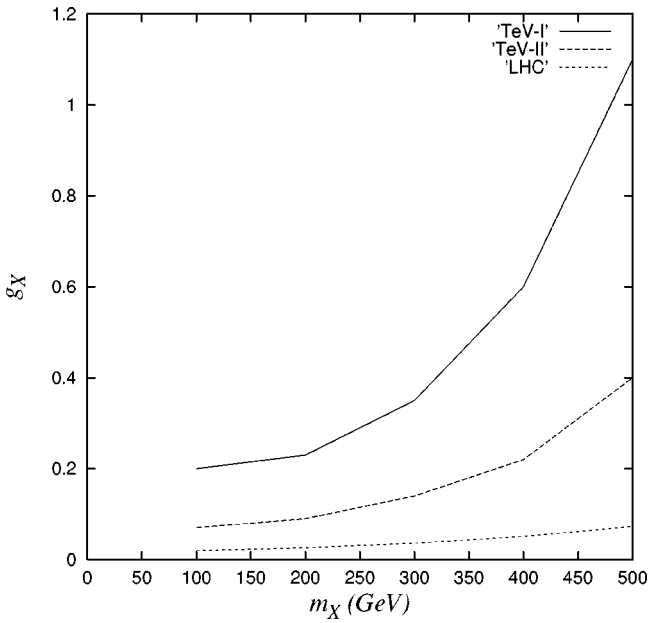

FIG. 2. The hadron collider limits on the X boson coupling as functions of its mass. The present limit from Tevatron (Run I) is shown along with the anticipated discovery limits for the Tevatron (Run II) and LHC.

presented their total $l\tau$ data from the Tevatron (run I), corresponding to an integrated luminosity of 110pb^{-1} [16]. The details of the data along with the selection cuts can be found in the second paper of [16]. A large fraction of this $l\tau$ data set contains a single accompanying jet. Of the 22 observed events in this data sample, 11 are estimated to come from

$$Z \to \tau\tau, \tag{23}$$

while most of the rest are estimated as τ fakes. We have estimated the X contribution to this channel using a parton level Monte Carlo program, which was found to reproduce the size of the above Z contribution reasonably well. As in the case of Z, the relevant production processes for X are the next leading order (NLO) Drell-Yan processes

$$q\bar{q} \to gX \quad \text{and} \quad gq(\bar{q}) \to q(\bar{q})X. \tag{24}$$

With 22 observed events and a background of similar magnitude, the 95% C.L. limit on the X boson signal can be estimated to be about 12 events [17]. The corresponding limit on g_X is shown as the solid line in Fig. 2. The plot is shown for $m_X \geqslant 100$ GeV, since for a light X the final lepton from $X \to \tau \to l$ decay becomes too soft to survive the selection cut. Thus there is a complementarity between the X search at hadron colliders and in Z-decay at LEP-I.

Note that the present Tevatron limit on g_X is not much better than the universality limit from LEP-I (Fig. 1). However, there is scope for significant improvement of this limit with much larger data samples expected from Tevatron (run II) and CERN Large Hadron Collider (LHC). In that case one can separate the X signal from Z by imposing a p_T cut on $X(Z)$, which will enable one to reconstruct the momenta of the decay τ pair and hence the $X(Z)$ mass. This technique has been widely investigated in the context of Higgs boson search at hadron colliders in the $\tau\tau$ channel. We have explored this quantitatively by imposing a $p_T^{X(Z)} > 50$ GeV cut for TeV-II and $p_T^{X(Z)} > 100$ GeV for the LHC. The resulting discovery limits of TeV-II and LHC are shown in Fig. 2. They correspond to 10 signal events with the expected luminosities of 2 fb^{-1} at TeV-II and 10 fb^{-1} at LHC. The latter corresponds to the low luminosity run of LHC. Even with this run it should be possible to probe for the X boson up to $m_X = 500$ GeV, assuming that g_X is of the same order of magnitude as g_1. The probe can be extended up to $m_X = 1$ TeV at the high luminosity run of LHC, which is expected to deliver an integrated luminosity of 100 fb^{-1}.

V. RADIATIVE NEUTRINO MASSES

In the presence of the $B\text{-}3L_\tau$ gauge symmetry, only ν_τ has a right-handed partner and thereby a seesaw mass. To accommodate the present data on neutrino oscillations, we need to allow ν_e and ν_μ to be massive, and have them mix with each other and ν_τ. To this end, we must break the remaining leptonic symmetries, i.e., multiplicative L_τ as well as additive L_e and L_μ. One possible scenario was already proposed in Ref. [4]. The scalar sector is extended to include a doublet

$$\begin{pmatrix} \eta^+ \\ \eta^0 \end{pmatrix} \sim (1,2,1/2;-3) \tag{25}$$

and a charged singlet

$$\chi^- \sim (1,1,-1;-3). \tag{26}$$

The doublet breaks L_e, L_μ, and L_τ separately but an overall multiplicative lepton number is preserved. It also generates flavor-changing couplings of X to the charged leptons, details of which will be discussed in the next section. The singlet $\nu_{\tau R}$ is now paired with one linear combination of the three left-handed neutrinos. It appears at first sight that there are then two massless neutrinos left. However, since the three lepton numbers are no longer individually conserved, these neutrinos necessarily pick up radiative masses. This generally happens in two loops through double W exchange [18], but the masses so obtained are extremely small. In the present scenario without the χ^- singlet, one of the two massless neutrinos at tree level does pick up a radiative mass in one loop [19], but it is also too small.

To obtain phenomenologically interesting radiative neutrino masses, we add the χ^- singlet to produce the following new interactions:

$$f_l(\nu_l \tau_L - l_L \nu_\tau)\chi^+, \quad (\phi^+ \eta^0 - \phi^0 \eta^+)\chi^- \chi^0, \tag{27}$$

where $l = e, \mu$. The mass-generating radiative mechanism of Ref. [20] is now operative, as shown in Fig. 3. One should note that the above scalar sector contains a pseudo-Goldstone boson which comes about because there are 3 global U(1) symmetries, corresponding each to rotating the phases of ϕ, η, and χ^0 independently in the Higgs potential, and only 2 local U(1) symmetries which get broken. How-

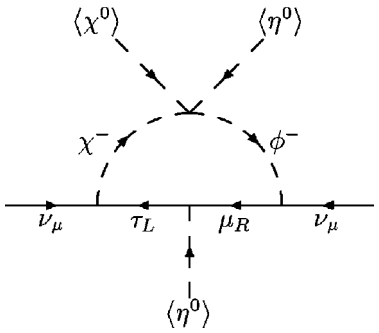

FIG. 3. Radiative mechanism for neutrino masses in the flavor-changing scenario.

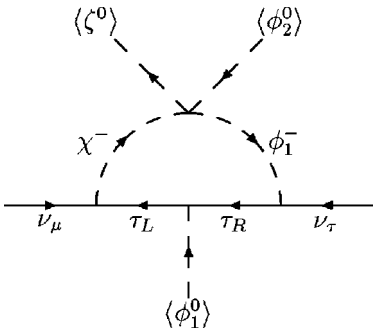

FIG. 4. Radiative mechanism for neutrino mass in the flavor-conserving scenario.

ever, if an extra neutral scalar ζ^0 transforming as (1,1,0; -3) is added, then the Higgs potential will have two more terms ($\eta^\dagger\phi\zeta^0$, $\zeta^0\zeta^0\chi^0$), and the extra unwanted U(1) symmetry is eliminated.

An alternative scenario is to replace η with a second Φ doublet, but retain χ^- as well as ζ^0. In that case, $\langle\phi^0_{1,2}\rangle$ break only $\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y$ whereas $\langle\chi^0\rangle$ and $\langle\zeta^0\rangle$ break only $\mathrm{U}(1)_X$. In contrast, in the other scenario, $\langle\eta^0\rangle$ breaks both. Hence X has no tree-level flavor-changing couplings, and does not mix with Z except through the cross kinetic-energy terms [21] which we assume to be negligible. We show in Fig. 4 the one-loop diagram connecting ν_μ with ν_τ. Note that in this scenario, only one linear combination of ν_e and ν_μ picks up a nonzero mass in one loop. The other linear combination will get a mass in two loops [18]. This structure is consistent with the new Super Kamiokande data. To suppress flavor-changing neutral-current interactions in the scalar sector, we impose a discrete Z_2 symmetry such that Φ_1 is even and Φ_2 is odd so that the latter does not couple to leptons. This discrete symmetry is then broken softly by the $\Phi_1^\dagger\Phi_2+\mathrm{H.c.}$ term in the Higgs potential, as in the minimal supersymmetric standard model.

VI. LEPTON AND QUARK FLAVOR NONCONSERVATION

In the scenario where we add the scalar $\eta\sim(1,2,1/2;-3)$ doublet, the charged-lepton mass matrix linking $\overline{e_L}$, $\overline{\mu_L}$, $\overline{\tau_L}$ to e_R, μ_R, τ_R can be chosen to be of the form

$$\mathcal{M}_l=\begin{bmatrix} m_e & 0 & 0 \\ 0 & m_\mu & 0 \\ a_e & a_\mu & m_\tau \end{bmatrix}. \tag{28}$$

Since X couples only to τ before symmetry breaking, the fact that $\mathcal{M}_l$ is not diagonal induces flavor-changing couplings of X to e and μ as follows:

$$3g_X X_\nu\left(\frac{a_\mu}{m_\tau}\overline{\mu_R}\gamma^\nu\tau_R+\frac{a_e}{m_\tau}\overline{e_R}\gamma^\nu\tau_R-\frac{a_\mu a_e}{m_\tau^2}\overline{e_R}\gamma^\nu\mu_R+\mathrm{H.c.}\right). \tag{29}$$

The best individual bounds on a_μ and a_e come from the nonobservation of $\tau\to\mu\pi^+\pi^-$ and $\tau\to e\pi^+\pi^-$. Assuming that the ratios of the above rates to that of $\tau\to\nu_\tau\pi^-\pi^0$ are roughly given by those of their inclusive rates, and using the upper limits [22] of 7.4×10^{-6} and 4.4×10^{-6} on their branching fractions, we find

$$\frac{g_X^2}{m_X^2}a_\mu m_\tau<3.8\times10^{-7},\quad \frac{g_X^2}{m_X^2}a_e m_\tau<2.9\times10^{-7}. \tag{30}$$

The best bound on the product $a_\mu a_e$ comes from the nonobservation of $\mu-e$ conversion in nuclei. Using the formalism of Ref. [23] and the experimental upper limit [22] of 4.3×10^{-12} for $\sigma(\mu\mathrm{Ti}\to e\mathrm{Ti})/\sigma(\mu\mathrm{Ti}\to\mathrm{capture})$, we find

$$\frac{g_X^2}{m_X^2}a_\mu a_e<3.1\times10^{-12}. \tag{31}$$

For $g_X=0.2$ and $m_X=60$ GeV, the above bounds translate to

$$a_\mu<19\ \mathrm{MeV},\quad a_e<15\ \mathrm{MeV},\quad a_\mu a_e<0.3\ (\mathrm{MeV})^2. \tag{32}$$

We note that [4] a radiative ν_μ mass of 2.3×10^{-3} eV could be obtained with $a_\mu=10$ MeV.

Independent of possible tree-level lepton flavor nonconservation in any variation of the B-$3L_\tau$ model, there is quark flavor nonconservation in one-loop order involving the X gauge boson. This contributes to decays such as $K^+\to\pi^+\nu_\tau\overline{\nu_\tau}$ and $b\to s\tau^+\tau^-$. However, since X has only vector couplings to quarks, the effective one-loop transition amplitude $q_1\to q_2X$ has the same form [24] as $q_1\to q_2\gamma$, i.e., $k^\mu\epsilon_X^\nu\overline{q_2}\sigma_{\mu\nu}(A+B\gamma_5)q_1$, where A and B are proportional to m_q/M_W^2. In contrast, the transition amplitude $q_1\to q_2Z$ is of the form [25] $\epsilon_Z^\nu\overline{q_2}\gamma_\nu(1-\gamma_5)q_1$. Hence the contribution of X to these amplitudes is suppressed by m_q^2/m_X^2 relative to that of Z and is always negligible.

VII. CONCLUSION

Since B-$3L_\tau$ can be gauged with the addition of $\nu_{\tau R}$, the possible existence of the associated gauge boson X should be investigated. We find its coupling g_X and mass m_X to be constrained by the data on Z decay in two important ways. For $m_X<70$ GeV, the nonobservation of the direct decay of Z to X gives an upper limit on g_X as shown in Fig. 1. For

$m_X < 50$ GeV, we find a rather stringent limit of $g_X < 0.1$. For $m_X > 56$ GeV, a better limit is obtained from the observed e-μ-τ universality of Z decay as shown also in Fig. 1.

The X boson may be produced at hadron colliders and be detected through its decay into τ pairs. The nonobservation of such events at the Tevatron puts an upper limit on g_X for $m_X \gtrsim M_Z$ as shown in Fig. 2. We have also estimated the discovery limits of X at the future run II of the Tevatron and at the LHC as shown also in Fig. 2. The latter offers a viable X boson signal up to $m_X \sim 1$ TeV.

Even though only ν_τ gets a tree-level mass in the B-$3L_\tau$ gauge model, the other two neutrinos may also acquire masses and mix with ν_τ through radiative corrections with an extended scalar sector. We have presented two possible scenarios, one with tree-level flavor-nondiagonal couplings of the X boson to charged leptons and one without. In the first scenario, important phenomenological constraints come from τ decay and μ-e conversion in nuclei. On the other hand, quark flavor nonconservation is always suppressed in one loop because X couples only vectorially.

In conclusion, if the X boson exists, it may be hiding very effectively even if its coupling is not too small and its mass is below M_Z. However, future hadron colliders can probe for the X boson up to a mass range of ~ 1 TeV if its coupling is of the same order of magnitude as g_1.

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