Mirror-lepton phenomenology in a left-right model with ultralight Dirac neutrinos

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A study of the production and decay of heavy mirror leptons is presented in the framework of a left-right model that yields ultralight Dirac neutrinos. Qualitative limits on the masses of the light neutrinos are also presented in terms of bounds on the heavy-mirror-lepton masses.

I. INGREDIENTS OF THE MODEL

In a recent paper[1] the authors have proposed a mechanism which generates ultralight Dirac neutrinos. Unlike other mechanisms[2] involving Majorana or Dirac neutrinos, no high-energy mass scales ($>10^9$ GeV) are involved. Instead, symmetry breaking at the electroweak scale is responsible for the mass generation. The framework of our calculation is a left-right-symmetric model containing mirror leptons.[3] The discrete symmetry[4] required to prevent the standard and mirror particles from pairing off and becoming massive can lead to a skewed form of the mass matrix, which results in some neutrinos acquiring ultralight masses (<100 eV). The mirror neutrinos, on the other hand, acquire masses in the 10—200-GeV range. The purpose of this paper is to pursue some of the phenomenology associated with these "nearby" mirror objects. Other recent related phenomenological studies of heavy leptons have been made by Gronau, Leung, and Rosner,[5] Maalampi and Mursula,[6] and Bagger, Dimopoulos, Massó, and Reno.[7]

We now summarize the main ingredients of our model. The weak gauge group is $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} = G_W$. The leptons are placed in doublets of $G_W \otimes Z$ as follows, where Z is the discrete symmetry (cf. Ref. 1):

$$\begin{aligned}
L_i &= \begin{bmatrix} \nu_i' \\ l_i^- \end{bmatrix}_L \sim (2,1;-1)_{X(L_i)} ,\\
L_i^c &= \begin{bmatrix} l_i^c \\ -\nu_i'^c \end{bmatrix}_L \sim (1,2;1)_{X(L_i^c)} ,\\
M_i &= \begin{bmatrix} E_i^+ \\ N_i' \end{bmatrix}_L \sim (2,1;1)_{X(M_i)} ,\\
M_i^c &= \begin{bmatrix} N_i'^c \\ -E_i^c \end{bmatrix}_L \sim (1,2;-1)_{X(M_i^c)} ,
\end{aligned} \tag{1.1}$$

in terms of left-handed two-component Weyl spinors. The L_i are the standard left-handed leptons observed in nature, and the L_i^c are their right-handed counterparts, which undergo right-handed weak interactions that have yet to be detected since the corresponding weak gauge fields W' and Z' are presumably much more massive than the observed W and Z. The mirror objects M_i, on the other hand, have the usual left-handed weak interactions but are presumably too heavy to have been observed so far.

Masses are generated by Higgs multiplets:

$$\begin{aligned}
\phi_i &= \begin{bmatrix} \phi^0 & \phi^+ \\ \phi^- & \phi'^0 \end{bmatrix}_i \sim (2,2;0)_{X(\phi_i)} ,\\
\Phi_i &= \begin{bmatrix} \Phi^+ & \Phi^{++} \\ \Phi^0 & \Phi'^+ \end{bmatrix}_i \sim (2,2;2)_{X(\Phi_i)} .
\end{aligned} \tag{1.2}$$

The ϕ and Φ fields transform as doublets under both $SU(2)_L$ and $SU(2)_R$ and so generate Dirac masses. No Majorana mass contributions are generated, even when radiative corrections are taken into account. The $SU(2)_R$ symmetry is broken at a scale greater than the electroweak scale by introducing the two Higgs fields

$$\chi_L = \begin{bmatrix} \chi^+ \\ \chi^0 \end{bmatrix} \sim (2,1;1), \quad \chi_R = \begin{bmatrix} \bar{\chi}^+ \\ \bar{\chi}^0 \end{bmatrix} \sim (1,2;1) \tag{1.3}$$

with the usual assumption that $\langle \chi_R \rangle \gg \langle \chi_L \rangle$.

Proper choice of the discrete-symmetry labels can restrict the Yukawa Lagrangian to the skewed form

$$\begin{aligned}
L_Y = {} & g_{LL} L^T(-i\sigma_2)\phi L^c + \tilde{g}_{MM} M^T(-i\sigma_2)\tilde{\phi} M^c \\
& + \tilde{g}_{ML} M^T(-i\sigma_2)\tilde{\Phi} L^c + \text{H.c.} ,
\end{aligned} \tag{1.4}$$

where the ϕ, $\tilde{\phi} = \sigma\phi^*\sigma_2$, and $\tilde{\Phi} = \sigma_2\Phi^*\sigma_2$ terms are all segregated. [This segregation also avoids flavor-changing neutral currents (FCNC's) mediated by Higgs bosons.] This is illustrated explicitly in Ref. 1 with a discrete Z_6 symmetry in the case of two standard and two mirror families. Mixing terms of the type

$$\mathrm{Tr}(\Phi_i^\dagger \Phi_j \phi_k^\dagger \tilde{\phi}_l), \quad i+j+k+l = 0 \ (\mathrm{mod}\ 6) \tag{1.5}$$

appearing in the Higgs Lagrangian violate a U(1) symmetry of the ϕ fields and permit a naturally small ratio of the vacuum expectation values

$$x = \langle \phi'^0 \rangle / \langle \phi^0 \rangle \ll 1 . \tag{1.6}$$

Because of the skewed segregated form of the Yukawa Lagrangian the neutral-lepton masses can be a factor of x^2 smaller than the charged-lepton masses and ultralight. In the next section we explore the mass matrices in some detail.

II. RESULTS FOR NO FAMILY MIXING

The charged-lepton and neutrino mass matrices have the following form in the basis $\{M,L,M^c,L^c\}$:

$$M_{\pm,0}=\begin{bmatrix}0 & M\\ M^T & 0\end{bmatrix}_{\pm,0}, \tag{2.1a}$$

where

$$M_{\pm}=\begin{array}{c|cc} & E^c & l^c\\ \hline E & A & 0\\ l & 0 & S\end{array}\;{}_{\pm}\,,$$

$$M_0=\begin{array}{c|cc} & N'^c & \nu'^c\\ \hline N' & A & B\\ \nu' & 0 & S\end{array}\;{}_{0}\,, \tag{2.1b}$$

and

$$A_{\pm}=\tilde{g}_{MM}\langle\phi^{0*}\rangle,\quad A_0=\tilde{g}_{MM}\langle\phi'^{0*}\rangle\,,$$
$$S_{\pm}=g_{LL}\langle\phi^{0}\rangle,\quad S_0=g_{LL}\langle\phi'^{0}\rangle\,, \tag{2.1c}$$
$$B_0=\tilde{g}_{ML}\langle\Phi^{0*}\rangle\,.$$

Each entry (A,S,B) is a matrix in the family space with A $N_m\times N_m$ dimensional, B $N_m\times N_s$ dimensional, and S $N_s\times N_s$ dimensional, where N_m is the number of mirror families and N_s the number of standard families. The skewed form of M_0 is a direct consequence of the discrete symmetry yielding Eq. (1.4).

In the following we shall restrict our attention to the simplified situation where no family mixing of the leptons occurs.[8] Hence A and S are diagonal matrices while B is diagonal only if $N_m=N_s$; i.e., the number of mirror and standard families are equal. We can then reorder the basis so that the M matrices are block family diagonal. For each family where the standard and mirror neutrinos are paired, the neutral mass matrix is

$$\begin{bmatrix}0 & 0 & a & b\\ 0 & 0 & 0 & s\\ a & 0 & 0 & 0\\ b & s & 0 & 0\end{bmatrix}_0\,. \tag{2.2}$$

From Eqs. (2.1), we see the neutrino entries a_0 and s_0 are scaled relative to the charged-lepton entries $a_{\pm}$ and $s_{\pm}$ by a factor of x in (1.6) which we take to be of order 10^{-3} consistent with a radiative origin:

$$x=a_0/a_{\pm}=s_0/s_{\pm}\sim 10^{-3}\,. \tag{2.3}$$

Here $a_{\pm}$ and $s_{\pm}$ are, respectively, just the mirror and standard charged-lepton masses, and thus, $a_{\pm,0}>s_{\pm,0}$. The neutral standard-mirror mixing entry b_0 is expected to be of the order of the electroweak scale or less in our model, while a_0 and s_0 are suppressed by a factor x. Hence, the expected ordering of terms is assumed to be $s\ll a\ll b$, where we have suppressed the subscript zero here and in the following, when no ambiguity arises. The neutral mass eigenvalues of (2.2) are then computed to be

$$\lambda\simeq\pm b(1+a^2/2b^2)\simeq\pm b\sim 100\text{ GeV}\,,$$
$$\lambda'\simeq\pm(as/b)(1-a^2/2b^2)\simeq\pm as/b\sim 100\text{ eV}\,. \tag{2.4}$$

For unpaired standard or mirror families, the neutral masses are given by s or a, respectively.

The corresponding eigenstates represent pairs of degenerate Majorana particles. From these we can construct massive Dirac particles with masses and eigenstates as follows:

$M=b(1+a^2/2b^2)\simeq b$:

$$N_L=N'_L\cos\beta+\nu'_L\sin\beta\simeq N'_L+(s/b)\nu'_L\,,$$
$$N^c_L=N'^c_L\sin\alpha+\nu'^c_L\cos\alpha\simeq(a/b)N'^c_L+\nu'^c_L\,; \tag{2.5}$$

$\mu=(as/b)(1-a^2/2b^2)\simeq as/b$:

$$\nu_L=\nu'_L\cos\beta-N'_L\sin\beta\simeq\nu'_L-(s/b)N'_L\,,$$
$$\nu^c_L=-\nu'^c_L\sin\alpha+N'^c_L\cos\alpha\simeq-(a/b)\nu'^c_L+N'^c_L\,, \tag{2.6}$$

where the mass eigenstates are written in terms of the weak eigenstates for paired families. For unpaired families the mass and weak eigenstates are the same.

The weak doublets in (1.1) can then be reexpressed in terms of the mass eigenstates according to

$$\begin{bmatrix}\nu'\\ l^-\end{bmatrix}_L\simeq\begin{bmatrix}\nu+(s/b)N\\ 0\end{bmatrix}_L,$$
$$\begin{bmatrix}l^c\\ -\nu'^c\end{bmatrix}_L\simeq\begin{bmatrix}l^c\\ (a/b)\nu^c-N^c\end{bmatrix}_L,$$
$$\begin{bmatrix}E^+\\ N'\end{bmatrix}_L\simeq\begin{bmatrix}E^+\\ -(s/b)\nu+N\end{bmatrix}_L, \tag{2.7}$$
$$\begin{bmatrix}N'^c\\ -E^c\end{bmatrix}_L\simeq\begin{bmatrix}\nu^c+(a/b)N^c\\ -E^c\end{bmatrix}_L$$

for paired families. The mirror mixing in the $SU(2)_L$ doublets is very small, $\sim s/b\sim 10^{-8}-10^{-6}$, while the standard-mirror family mixing in the $SU(2)_R$ doublets though small is larger, $\sim a/b\sim 10^{-4}-10^{-3}$. Note that the conjugate doublets pair light with heavy fields to first approximation.

To first order for the charged W's, but only to zeroth order for the neutral Z's, the physical gauge-boson masses are determined to be

$$M_W\simeq g_L(\omega+\Omega-\omega'^2/\eta_R)^{1/2}\,,$$
$$M'_W\simeq g_R(\eta_R+\omega+\Omega)^{1/2}\,,$$
$$M_Z\simeq(g/\cos\theta_W)(\omega+\Omega)^{1/2}\,, \tag{2.8}$$
$$M'_Z\simeq(g'\cot\theta_W)(\eta_R+2\Omega)^{1/2}\,,$$

where from now on we set $g_L=g_R=g$, so that the coupling parameters in the left-right-symmetric model are related to those of the Weinberg-Salam model by

$$\sin\theta_W=g'/(g^2+2g'^2)^{1/2}\,,$$
$$g\sin\theta_W=g'\cos^{1/2}2\theta_W=e\,. \tag{2.9}$$

As a shorthand notation we have set

$$\eta_{L,R}=\langle\chi^0_{L,R}\rangle^2,\ \ \Omega=\langle\Phi^0\rangle^2\ ,$$
$$\omega=\langle\phi^0\rangle^2+\langle\phi'^0\rangle^2,\ \ \omega'=\langle\phi^0\rangle\langle\phi'^0\rangle\ , \tag{2.10}$$

and neglect η_L. In the same approximation the mass eigenstates are given by

$$\begin{aligned}
W&=W_L\cos\gamma+W_R\sin\gamma\simeq W_L+(\omega'/\eta_R)W_R\ ,\\
W'&=-W_L\sin\gamma+W_R\cos\gamma\simeq-(\omega'/\eta_R)W_L+W_R\ ,\\
Z&=\cos\theta_W[W^3_L-g'^2/(g^2+g'^2)W^3_R-gg'/(g^2+g'^2)B]\ ,\\
Z'&=g/(g^2+g'^2)^{1/2}[W^3_R-(g'/g)B]\ ,\\
A&=\sin\theta_W[W^3_L+W^3_R+(g'/g)B]\ .
\end{aligned} \tag{2.11}$$

The left-right mixing parameter for the charged W's is small,

$$\begin{aligned}
\omega'/\eta_R&=\langle\phi^0\rangle\langle\phi'^0\rangle/\langle\chi^0_R\rangle^2\\
&\le10^{-3}(10^2)^2/(10^4)^2\sim10^{-7}\ ,
\end{aligned} \tag{2.12}$$

well below the experimental limit[9] of 5×10^{-3}. Higher-order terms also introduce a mixing between the Z and Z' of order $(\omega+\Omega)/\eta_R$.

By combining Eqs. (2.7) and (2.11) we then obtain the physical gauge couplings:

$$\begin{aligned}
&(g/\sqrt{2})[\bar\nu'_L\gamma_\mu l_L+\bar E_L\gamma_\mu N'_L+(\omega'/\eta_R)(\bar\nu'_R\gamma_\mu l_R+\bar E_R\gamma_\mu N'_R)]W^+_\mu+\text{H.c.}\\
&\quad+(g/\sqrt{2})[\bar\nu'_R\gamma_\mu l_R+\bar E_R\gamma_\mu N'_R-(\omega'/\eta_R)(\bar\nu'_L\gamma_\mu l_L+\bar E_L\gamma_\mu N'_L)]W'^+_\mu+\text{H.c.}\\
&\quad+(g/2\cos\theta_W)[\bar\nu'_L\gamma_\mu\nu'_L-\bar l_L\gamma_\mu l_L-\bar N'_L\gamma_\mu N'_L+\bar E_L\gamma_\mu E_L-2\sin^2\theta_W(-\bar l\gamma_\mu l+\bar E\gamma_\mu E)]Z_\mu\\
&\quad+(g'/2)\tan\theta_W[\bar\nu'_L\gamma_\mu\nu'_L-\bar N'_L\gamma_\mu N'_L+\bar\nu'_R\gamma_\mu\nu'_R-\bar N'_R\gamma_\mu N'_R\\
&\qquad+(\cos2\theta_W/\sin^2\theta_W)(\bar\nu'_R\gamma_\mu\nu'_R-\bar l_R\gamma_\mu l_R-\bar N'_R\gamma_\mu N'_R+\bar E_R\gamma_\mu E_R)-(-\bar l\gamma_\mu l+\bar E\gamma_\mu E)]Z'_\mu\\
&\quad+e(-\bar l\gamma_\mu l+\bar E\gamma_\mu E)A_\mu\ .
\end{aligned} \tag{2.13}$$

The physical W exhibits a small right-handed coupling to the leptons which is down by a factor of $\sim10^{-7}$ as noted in (2.12). This is well within the present experimental limit[10] of 10^{-5}. The zeroth-order couplings of the W and Z are exactly as in the standard $SU(2)_L\otimes U(1)_Y$ model aside from the extra mirror contributions and admixtures of mirror leptons that follow from (2.7).

We close this section by presenting mass estimates for the neutrino and charged-mirror leptons. Again we restrict our attention to the no-family-mixing scenario for simplicity. As input we make use of the present upper limits[11] on the neutrino masses

$$\begin{aligned}
&m_{\nu_e}\le42\text{ eV}\ ,\\
&m_{\nu_\mu}\le250\text{ keV}\ ,\\
&m_{\nu_\tau}\le70\text{ MeV}\ ,
\end{aligned} \tag{2.14a}$$

as well as restrictions on the heavy-mirror particles

$$\begin{aligned}
&25\le M_E\le250\text{ GeV}\ ,\\
&7\le M_N\le250\text{ GeV}\ ,
\end{aligned} \tag{2.14b}$$

where the lower limits come from experiments[12] at the SLAC and DESY storage rings PEP and PETRA and the upper limits eliminate the possibility of a strong-interaction-strength Yukawa coupling.

With $x=10^{-3}$ for all families as proposed in (2.3), consistent with a radiative origin for this ratio, from the mass eigenvalues in (2.4) we deduce the mass relation

$$m_\nu M_N\simeq x^2m_lM_E\simeq10^{-6}m_lM_E \tag{2.15}$$

which connects the neutrino and neutral-mirror-lepton mass to the charged standard and mirror-lepton masses. Applying this relation to all three families, we find the light neutrinos are bounded by

$$\begin{aligned}
&0.05\text{ eV}<m_{\nu_e}\simeq0.5\text{ eV}<18\text{ eV}\ ,\\
&10\text{ eV}<m_{\nu_\mu}\simeq100\text{ eV}<3.6\text{ keV}\ ,\\
&180\text{ eV}<m_{\nu_\tau}\simeq1.8\text{ keV}<64\text{ keV}\ .
\end{aligned} \tag{2.16}$$

The central values emerge if the mirror-lepton masses M_N and M_E are taken equal; these correspond to the suggestions[13] of Wolfenstein and Rosen that the neutrino masses may be scaled relative to the corresponding charged-lepton masses. The lower limits appear when $M_E=25$ GeV and $M_N=250$ GeV, and finally the upper limits correspond to $M_N=7$ GeV and $M_E=250$ GeV. However, these bounds are crude estimates only, since x is not known. If only two mirror families exist and the third standard family remains unpaired, we would expect, on the other hand, $m_{\nu_\tau}=1.8$ MeV.

Mass relation (2.15) along with (2.14) also results in lower bounds on the ratio of the mirror leptons:

$$\begin{aligned}
&M_{N_E}/M_E\ge1.2\times10^{-2}\ ,\\
&M_{N_M}/M_M\ge2.1\times10^{-4}\ ,\\
&M_{N_T}/M_T\ge2.5\times10^{-5}\ .
\end{aligned} \tag{2.17}$$

It is clear that these inequalities impose almost no restrictions on the mirror leptons, so we shall refer to allowed ranges in (2.14b) when investigating the mirror-lepton phenomenology.

III. PRODUCTION OF MIRROR LEPTONS

A. $W \to E\bar{N}'$ and $Z \to E\bar{E}, N'\bar{N}'$

1. Drell-Yan production of W and Z

If phase space permits, the mirror leptons can be produced singly or in pairs through decays of W and Z bosons. The Drell-Yan production process is well documented theoretically and has been experimentally observed[14] for the past three years at CERN. We refer the reader to Eichten, Hinchliffe, Lane, and Quigg,[15] for example, and do not repeat the formulas here.

Single-mirror-lepton production can occur through the small admixtures in Eq. (2.7), but the decay-rate formula

$$w(W^+ \to E^+ \bar{\nu}) = \frac{G M_W{}^3}{6\sqrt{2}\pi}\left[\frac{s}{b}\right]^2 \left[1+\frac{M_E{}^2}{2M_W{}^2}\right] \times \left[1-\frac{M_E{}^2}{M_W{}^2}\right]^2$$

$$\leq (10^{-16}-10^{-12})0.24 \text{ GeV} \qquad (3.1)$$

indicates a negligibly small branching ratio. On the other hand, double-mirror-lepton production, if kinematically allowed, is not so suppressed:

$$\begin{aligned}
w(W^+ \to E^+ \bar{N}) &= \frac{G M_W{}^3}{6\sqrt{2}\pi}\left[1-\frac{M_E{}^2+M_N{}^2}{2M_W{}^2}-\frac{(M_E{}^2-M_N{}^2)^2}{2M_W{}^4}\right]\left[1+\frac{M_E{}^2-M_N{}^2}{M_W{}^2}\right]\left[1-\frac{4M_W{}^2 M_E{}^2}{(M_W{}^2+M_E{}^2-M_N{}^2)^2}\right]^{1/2}, \\
w(Z^0 \to N\bar{N}) &= \frac{G M_Z{}^3}{12\sqrt{2}\pi}\left[1-\frac{M_N{}^2}{M_Z{}^2}\right]\left[1-4\frac{M_N{}^2}{M_Z{}^2}\right]^{1/2} \qquad (3.2) \\
w(Z^0 \to E\bar{E}) &= \frac{G M_Z{}^3}{12\sqrt{2}\pi}(1-4x_W+8x_W{}^2)\left[1-\frac{M_E{}^2}{M_Z{}^2}\right]\left[1-4\frac{M_E{}^2}{M_Z{}^2}\right]^{1/2},
\end{aligned}$$

where $x_W = \sin^2\theta_W$. Graphs of the W- and Z-boson partial decay rates as functions of the heavy-mirror masses are depicted in Figs. 1 and 2. From these figures it is clear that, if kinematically allowed, the mirror-lepton channels are sizable fractions of the light-lepton-channel decay rate, $w(W^+ \to l^+ \nu_l) \simeq 0.24$ GeV and $w(Z \to \nu\bar{\nu}) \simeq 0.18$ GeV.

2. Drell-Yan production of virtual W's and Z's

If the heavy leptons are too massive to be counted as real decay channels of the W and Z, their production can still occur through virtual W and Z contributions to the Drell-Yan process: $u\bar{d} \to W^+_{\text{virtual}} \to E\bar{N}$ and $q\bar{q} \to Z_{\text{virtual}} \to N\bar{N}, E\bar{E}$. The differential cross section for the former is given by

$$\frac{d\hat{\sigma}}{dt} = \frac{\pi\alpha^2}{12\sin^4\theta_W}\frac{1}{s^2}\frac{(u-M_E{}^2)(u-M_N{}^2)}{(s-M_W{}^2)^2+M_W{}^2\Gamma_W{}^2}\ . \qquad (3.3)$$

This can be integrated over the quark-parton distribution functions to yield the total Drell-Yan cross section. Unfortunately, the production cross section via this mechanism falls off rapidly above the Z mass, and heavy-lepton pair production in this region is best explored in e^+e^- annihilation at the future CERN LEP collider.

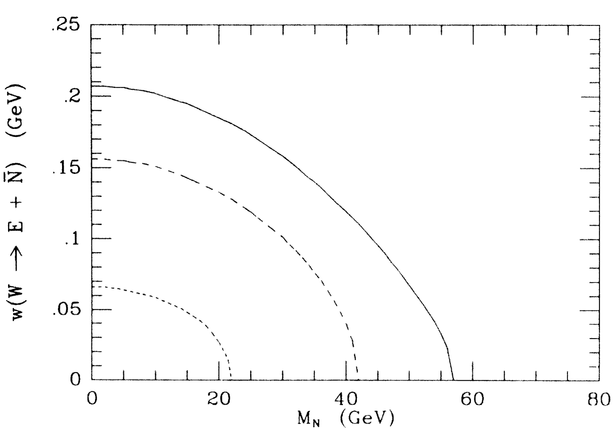

FIG. 1. Decay rate for $W \to E\bar{N}$ as a function of the mass of the N mirror lepton for various choices of the E mirror-lepton mass. From top to bottom the three curves correspond to $M_E = 25$, 40, and 60 GeV.

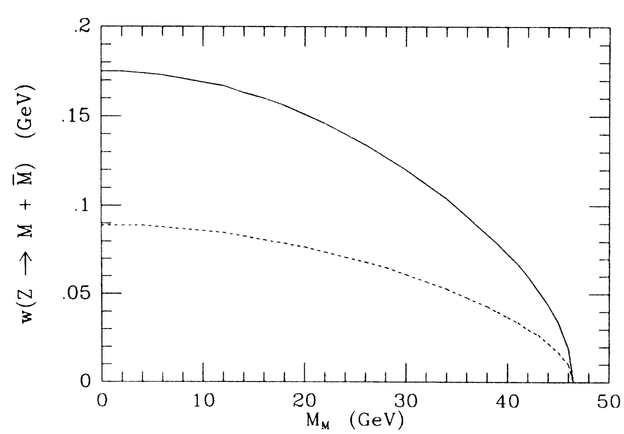

FIG. 2. Decay rate for Z into two mirror leptons. The solid (dashed) curve represents the $N\bar{N}$ ($E\bar{E}$) channel.

B. $W' \to E\bar{\nu}, N\bar{l}$ and $Z' \to N\bar{N}, E\bar{E}$

Drell-Yan production of the $SU(2)_R$ gauge bosons W' and Z' will yield one or two mirror leptons upon decay as seen from Eqs. (2.7) and (2.13). No suppression factors need enter in the decay amplitudes; however, the present lower limits[16] on the masses of these $SU(2)_R$ bosons are of the order of 5 TeV, so direct production cannot be achieved until the proposed Superconducting Super Collider is in operation. At center-of-mass energies accessible to the Fermilab Tevatron Collider, the virtual W' and Z' are too far off mass shell to yield measurable cross sections. We do not pursue this matter further here. It is clear that the model will stand or fall by the production processes discussed in Sec. III A.

IV. MIRROR-LEPTON DECAYS

Once the mirror leptons are produced, the heavier ones can decay into the lighter ones rapidly without suppression via the $SU(2)_L$ weak-boson channels. The lightest mirror, on the other hand, will be much longer lived and decays either through the virtual $SU(2)_R$ bosons or through the $SU(2)_L$ bosons, where the very small admixture of standard and mirror leptons can play a role. As indicated in Sec. II, the present constraints on the mirror spectrum are rather loose. For this reason, we consider both cases: $M_N > M_E$ and $M_E > M_N$.

The decay rate for the heavier mirror into the lighter mirror lepton is given by

$$w(M_1 \to M_2 l \nu_l) = \frac{G^2}{3(2\pi)/3} M_1 M_2{}^4 \times [(2\sec^3\theta - 5\sec\theta)\tan\theta + 3\ln|M_N/M_E|] , \qquad (4.1)$$

where

$$\tan\theta = \frac{M_1{}^2 - M_2{}^2}{2M_1M_2}, \quad \sec\theta = \frac{M_1{}^2 + M_2{}^2}{2M_1M_2}$$

for the purely leptonic channels. The results are graphed in Fig. 3 for different combinations of mirror-lepton masses. The decay rate for $M_1 \to M_2 q\bar{q}$ is three times larger due to the color factor, if light quarks are emitted.

If N is the lightest mirror lepton, its dominant decay occurs semileptonically through a virtual W' boson with rate given by

$$w(N \to l q\bar{q}') = \frac{3G^2}{192\pi^3} M_N{}^5 (M_W/M_{W'})^4 . \qquad (4.2)$$

Although a severe suppression arises through the present upper bound[16] $M_W/M_{W'} \le 0.02$, the only competing decay mode arises through the tiny admixture of standard and mirror leptons in the left-handed doublets of (2.7). The corresponding decay rate is given by

$$w(N \to l q\bar{q}') = \frac{3G^2}{192\pi^3} M_N{}^5 (s/b)^2 . \qquad (4.3)$$

The results for the (4.2) decay mode are presented in Fig.

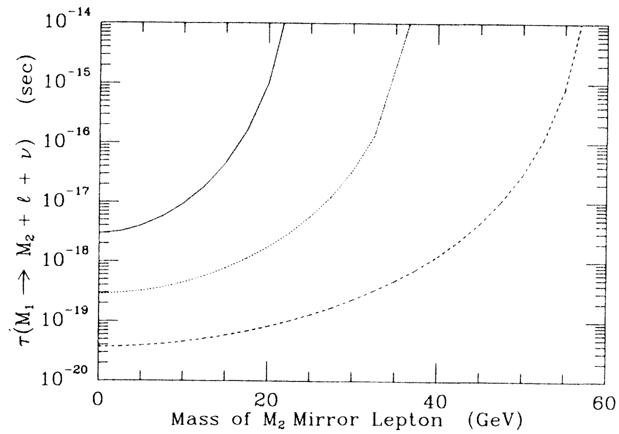

FIG. 3. Lifetime of a heavy mirror lepton M_1 decaying into a lighter mirror lepton M_2 as a function of the M_2 mass. From top to bottom the three curves correspond to an M_1 mass of 25, 40, and 60 GeV.

4 for different masses of the W' boson. A similar situation exists if E is the lightest mirror lepton. The dominant decay is then $E \to \nu q\bar{q}'$ with the analog of the decay rate (4.2) dominating that of (4.3) again by several orders of magnitude. Although the suppression factor in Eq. (4.2), for example, is very great, the lightest mirror lepton is not highly stable but is expected to decay with a typical weak-interaction lifetime. This makes the light-mirror decay especially accessible to experimental study.

V. LEPTON MIXING AND FCNC DECAYS

For a general model of the type we consider here there is the possibility of mixing between lepton families in the mass matrix. For left-handed neutrinos, mixing between light and heavy neutrinos continues to be suppressed by a factor $s/b \sim 10^{-6}$, so that the mixing of the light neutrinos among themselves is unitary to a good approximation.

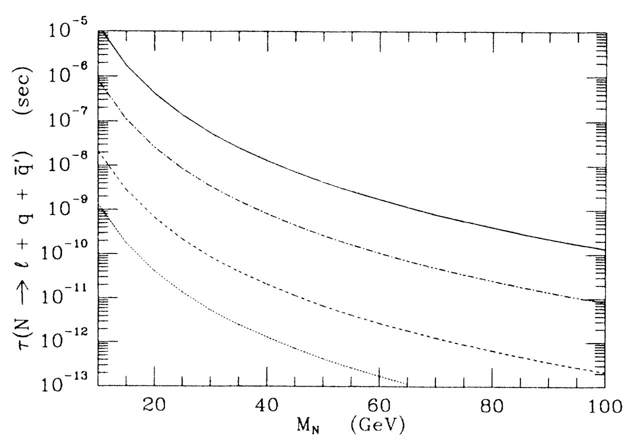

FIG. 4. Lifetime of the lightest mirror lepton decaying according to Eq. (4.2) as a function of its mass. From top to bottom the four curves correspond to a W' mass of 50, 25, 10, and 5 TeV.

Thus, the usual Z coupling is essentially flavor conserving and obeys universality. This unitarity also ensures that left-handed charged-current radiative contributions to neutral-current processes like $\mu\to e\gamma$ are suppressed by the usual Glashow-Iliopoulos-Maiani-type mechanism.

For right-handed neutrinos the mixing of the light neutrinos is unitary only up to corrections of order $a/b\sim 10^{-3,4}$, but the resulting departure from universality of the Z' coupling is insignificant due to the large mass of the Z'. Also the right-handed charged current connects light charged leptons to heavy rather than light neutral fermions. This allows the possibility of single or pair production of heavy neutral leptons in an ep or e^+e^- collision, though this again will be suppressed by the large mass of the W'. Furthermore, it implies that the right-handed contribution to, e.g., $\mu\to e\gamma$ is proportional to $(M_N{}^2/M_{W'}{}^2)^2$, where M_N is the mass of a heavy neutral, with no a/b suppression. Nevertheless, this suppression is still sufficient to avoid contradicting the current experimental limit[17] on the branching ratio of 1.7×10^{-10}.

There are also potential flavor-changing effects in charged-current processes. These are scaled by the leptonic equivalent of the Cabibbo or Kobayashi-Maskawa mixing angles, which by analogy we can assume to be small. For example, in the case of mixing between the first two lepton families, we find explicitly that the mixing of the left-handed ν_e and ν_μ is of order s'/m_μ, where s' is the element of the mass matrix that gives the mixing of the e^- and μ^-, and we have made use of the fact that $m_\mu \gg m_e$. The lepton-number violation can thus be safely small if s' is small in spite of the unusual mixing with the mirror neutrinos. As indicated above, right-handed charged-current neutrino scattering would involve the creation of a heavy particle and is suppressed for this reason as well as by the mass of W'.

Flavor-changing effects due to Higgs particles can be more significant. However, by including two χ_R fields, we can arrange to obtain just two light physical Higgs particles with masses on the order of the $SU(2)_L$-breaking scale. The rest have masses on the order of the scale Λ_R of $SU(2)_R$ symmetry breaking, and their interactions, flavor changing and otherwise, are suppressed. The two light Higgs bosons include one associated with the development of the $SU(2)_L$-breaking vacuum expectation value, which always has flavor-diagonal couplings to fermions, while the other couples to fermions only like ~ 100 GeV$/\Lambda_R$, and again its effects would not have been seen.

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