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Mirror Lepton Phenomenology in a
Left-Right Model with Ultralight Dirac Neutrinos

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Abstract

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:

$$
L_i = \begin{pmatrix} \nu_i' \\ \ell_i^- \end{pmatrix}_L \sim (2,1;-1)_{X(L_i)}, \qquad
L_i^c = \begin{pmatrix} \ell_i^c \\ -\nu_i'^c \end{pmatrix}_L \sim (1,2;1)_{X(L_i^c)},
$$

$$
M_i = \begin{pmatrix} E_i^+ \\ N_i' \end{pmatrix}_L \sim (2,1;1)_{X(M_i)}, \qquad
M_i^c = \begin{pmatrix} N_i'^c \\ -E_i^c \end{pmatrix}_L \sim (1,2;-1)_{X(M_i^c)}
\tag{1.1}
$$

in terms of lefthanded two-component Weyl spinors. The L_i are the standard lefthanded leptons observed in nature, and the L_i^c are their righthanded counterparts, which undergo righthanded 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 lefthanded weak interactions but are presumably too heavy to have been observed

so far.

Masses are generated by Higgs multiplets

$$
\phi_i = \begin{pmatrix} \phi^0 & \phi^+ \\ \phi^- & \phi'^0 \end{pmatrix}_i \sim (2,2;0)_{X(\phi_i)}, \qquad \Phi_i = \begin{pmatrix} \Phi^+ & \Phi^{++} \\ \Phi^0 & \Phi'^+ \end{pmatrix}_i \sim (2,2;2)_{X(\Phi_i)} \quad (1.2)
$$

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

$$
\chi_L = \begin{pmatrix} \chi^+ \\ \chi^0 \end{pmatrix} \sim (2,1;1), \qquad \chi_R = \begin{pmatrix} \overline{\chi}^+ \\ \overline{\chi}^0 \end{pmatrix} \sim (1,2;1) \qquad (1.3)
$$

with the usual assumption that $< \chi_R > \gg < \chi_L >$.

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

$$
L_Y = g_{LL} L^T (-i\sigma_2)\phi L^c + \tilde{g}_{MM} M^T (-i\sigma_2)\tilde{\phi} M^c + \tilde{g}_{ML} M^T (-i\sigma_2)\tilde{\Phi} L^c + h.c. \quad (1.4)
$$

where the ϕ, $\tilde{\phi} = \sigma_2 \phi^* \sigma_2$ and $\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2$ terms are all segregated. (This segregation also avoids flavor-changing neutral currents (FCNC) 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

$$
Tr(\Phi_i^\dagger \Phi_j \phi_k^\dagger \tilde{\phi}_\ell), \qquad i + j + k + \ell = 0 \qquad (mod\ 6) \qquad (1.5)
$$

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

$$
x = < \phi'^0 > / < \phi^0 > \ll 1 \qquad (1.6)
$$

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 $\{M, L, M^c, L^c\}$:

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

where

$$M_{\pm} = \begin{array}{c} \\ E \\ \ell \end{array} \begin{array}{cc} E^c & \ell^c \\ \begin{bmatrix} A & 0 \\ 0 & S \end{bmatrix}_{\pm} \end{array}$$

$$\tag{2.1b}$$

$$M_0 = \begin{array}{c} \\ N' \\ \nu' \end{array} \begin{array}{cc} N'^c & \nu'^c \\ \begin{bmatrix} A & B \\ 0 & S \end{bmatrix}_0 \end{array}$$

and

$$A_{\pm} = \tilde{g}_{MM} < \phi^{0*} >, \qquad A_0 = \tilde{g}_{MM} < \phi'^{0*} >,$$

$$S_{\pm} = g_{LL} < \phi^0 >, \qquad S_0 = g_{LL} < \phi'^0 >, \tag{2.1c}$$

$$B_0 = \tilde{g}_{ML} < \Phi^{0*} >$$

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

$$\begin{pmatrix} 0 & 0 & a & b \\ 0 & 0 & 0 & s \\ a & 0 & 0 & 0 \\ b & s & 0 & 0 \end{pmatrix}_0 \tag{2.2}$$

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:

$$x = a_0/a_\pm = s_0/s_\pm \sim O(10^{-3}) \tag{2.3}$$

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

$$\lambda \simeq \pm b(1 + a^2/2b^2) \simeq \pm b \sim 100 GeV$$

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

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:

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

$$N_L = N_L' \cos\beta + \nu_L' \sin\beta \simeq N_L' + (s/b)\nu_L'$$

$$N_L^c = N_L'^c \sin\alpha + \nu_L'^c \cos\alpha \simeq (a/b)N_L'^c + \nu_L'^c \qquad (2.5)$$

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

$$\nu_L = \nu_L' \cos\beta - N_L' \sin\beta \simeq \nu_L' - (s/b)N_L'$$

$$\nu_L^c = -\nu_L'^c \sin\alpha + N_L'^c \cos\alpha \simeq -(a/b)\nu_L'^c + N_L'^c \qquad (2.6)$$

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

$$\begin{pmatrix} \nu' \\ \ell^- \end{pmatrix}_L \simeq \begin{pmatrix} \nu + (s/b)N \\ \ell^- \end{pmatrix}_L, \qquad \begin{pmatrix} \ell^c \\ -\nu'^c \end{pmatrix}_L \simeq \begin{pmatrix} \ell^c \\ (a/b)\nu^c - N^c \end{pmatrix}_L,$$

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

for paired families. The mirror mixing in the $SU(2)_L$ doublets is very small, $O(s/b) \sim 10^{-8} - 10^{-6}$, while the standard-mirror family mixing in the $SU(2)_R$ doublets though small is larger, $O(a/b) \sim 10^{-4} - 10^{-3}$. 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

$$M_W \simeq g_L(\omega + \Omega - \omega'^2/\eta_R)^{1/2}$$

$$M_W' \simeq g_R(\eta_R + \omega + \Omega)^{1/2}$$

$$M_Z \simeq (g/\cos\theta_W)(\omega + \Omega)^{1/2} \qquad (2.8)$$

$$M_Z' \simeq (g'\cot\theta_W)(\eta_R + 2\Omega)^{1/2}$$

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

$$\sin\theta_W = g'/\sqrt{g^2 + 2g'^2},$$

$$g\sin\theta_W = g'\cos^{1/2}2\theta_W = e \qquad\qquad (2.9)$$

As a shorthand notation we have set

$$\eta_{L,R} = <\chi^0_{L,R}>^2, \qquad\qquad \Omega = <\Phi^0>^2$$

$$\omega = <\phi^0>^2 + <\phi'^0>^2, \qquad \omega' = <\phi^0><\phi'^0> \qquad (2.10)$$

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

$$W = W_L\cos\gamma + W_R\sin\gamma \simeq W_L + (\omega'/\eta_R)W_R$$

$$W' = -W_L\sin\gamma + W_R\cos\gamma \simeq -(\omega'/\eta_R)W_L + W_R$$

$$Z = \cos\theta_W\left[W^3_L - g'^2/(g^2+g'^2)W^3_R - gg'/(g^2+g'^2)B\right] \qquad (2.11)$$

$$Z' = g/\sqrt{g^2+g'^2}\left[W^3_R - (g'/g)B\right]$$

$$A = \sin\theta_W\left[W^3_L + W^3_R + (g'/g)B\right]$$

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

$$\omega'/\eta_R = <\phi^0><\phi'^0>/<\chi^0_R>^2 \leq 10^{-3}(10^2)^2/(10^4)^2 \sim 10^{-7} \qquad (2.12)$$

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:

$$(g/\sqrt{2})\Big[\bar{\nu}'_L\gamma_\mu\ell_L + \bar{E}_L\gamma_\mu N'_L + (\omega'/\eta_R)(\bar{\nu}'_R\gamma_\mu\ell_R + \bar{E}_R\gamma_\mu N'_R)\Big]W^+_\mu + h.c.$$

$$+ (g/\sqrt{2})\Big[\bar{\nu}'_R\gamma_\mu\ell_R + \bar{E}_R\gamma_\mu N'_R - (\omega'/\eta_R)(\bar{\nu}'_L\gamma_\mu\ell_L + \bar{E}_L\gamma_\mu N'_L)\Big]W'^+_\mu + h.c.$$

$$+ (g/2\cos\theta_W)\Big[\bar{\nu}'_L\gamma_\mu\nu'_L - \bar{\ell}_L\gamma_\mu\ell_L - \bar{N}'_L\gamma_\mu N'_L + \bar{E}_L\gamma_\mu E_L$$

$$- 2\sin^2\theta_W\left(-\bar{\ell}\gamma_\mu\ell + \bar{E}\gamma_\mu E\right)\Big]Z_\mu \tag{2.13}$$

$$+ (g'/2)\tan\theta_W\Big[\bar{\nu}'_L\gamma_\mu\nu'_L - \bar{N}'_L\gamma_\mu N'_L + \bar{\nu}'_R\gamma_\mu\nu'_R - \bar{N}'_R\gamma_\mu N'_R$$

$$+ (\cos 2\theta_W/\sin^2\theta_W)(\bar{\nu}'_R\gamma_\mu\nu'_R - \bar{\ell}_R\gamma_\mu\ell_R - \bar{N}'_R\gamma_\mu N'_R$$

$$+ \bar{E}_R\gamma_\mu E_R) - (-\bar{\ell}\gamma_\mu\ell + \bar{E}\gamma_\mu E)\Big]Z'_\mu$$

$$+ e(-\bar{\ell}\gamma_\mu\ell + \bar{E}\gamma_\mu E)A_\mu$$

The physical W exhibits a small righthanded coupling to the leptons which is down by a factor of $O(10^{-7})$ 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

$$m_{\nu_e} \leq 42\ eV$$

$$m_{\nu_\mu} \leq 250\ keV \tag{2.14a}$$

$$m_{\nu_\tau} \leq 70\ MeV$$

as well as restrictions on the heavy mirror particles

$$25GeV \leq M_E \leq 250GeV$$

$$7GeV \leq M_N \leq 250GeV$$

(2.14b)

where the lower limits come from experiments[12] at 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

$$m_\nu M_N \simeq x^2 m_\ell M_E \simeq 10^{-6} m_\ell M_E$$

(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

$$0.05eV < m_{\nu_e} \simeq 0.5eV < 18eV$$

$$10eV < m_{\nu_\mu} \simeq 100eV < 3.6keV$$

(2.16)

$$180eV < m_{\nu_\tau} \simeq 1.8keV < 64keV$$

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 $M_E = 25GeV$ and $M_N = 250GeV$, 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 2 mirror families exist and the third standard family remains unpaired, we would expect on the other hand, $m_{\nu_\tau} = 1.8MeV$.

Mass relation (2.15) along with (2.14) also result in lower bounds on the ratio

of the mirror leptons:

$$M_{N_E}/M_E \geq 1.2 \times 10^{-2}$$

$$M_{N_M}/M_M \geq 2.1 \times 10^{-4} \qquad\qquad (2.17)$$

$$M_{N_T}/M_T \geq 2.5 \times 10^{-5}$$

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, Hinchcliff, Lane and Quigg,[15] for example, and do not repeat the formulas here.

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

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

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

indicates a negligibly small branching ratio. On the other hand, double mirror

production, if kinematically allowed, is not so suppressed:

$$w(W^+ \to E^+ \bar{N}) = \frac{GM_W^3}{6\sqrt{2}\pi}\left[1 - \frac{M_E^2 + M_N^2}{2M_W^2} - \frac{(M_E^2 - M_N^2)^2}{2M_W^4}\right]$$

$$\cdot \left(1 + \frac{M_E^2 - M_N^2}{M_W^2}\right)\sqrt{1 - \frac{4M_W^2 M_E^2}{(M_W^2 + M_E^2 - M_N^2)^2}}$$

$$\tag{3.2}$$

$$w(Z^0 \to N\bar{N}) = \frac{GM_Z^3}{12\sqrt{2}\pi}\left(1 - \frac{M_N^2}{M_Z^2}\right)\sqrt{1 - 4\frac{M_N^2}{M_Z^2}}$$

$$w(Z^0 \to E\bar{E}) = \frac{GM_Z^3}{12\sqrt{2}\pi}(1 - 4x_W + 8x_W^2)\left(1 - \frac{M_E^2}{M_Z^2}\right)\sqrt{1 - 4\frac{M_E^2}{M_Z^2}}$$

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, $w(W^+ \to \ell^+ \nu_\ell) \simeq 0.24$ 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: $u\bar{d} \to W^+_{virtual} \to E\bar{N}$ and $q\bar{q} \to Z_{virtual} \to N\bar{N}, E\bar{E}$. The differential cross section for the former is given by

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

$$\tag{3.3}$$

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 LEP collider.

B. $W' \to E\bar{\nu}, N\bar{\ell}$ 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 can not be achieved until the proposed SSC 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 part A of this Section.

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 Section 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

$$w(M_1 \to M_2 \ell \nu_\ell) = \frac{G^2}{3(2\pi)/3} M_1 M_2^4 \left[(2\sec^3\theta - 5\sec\theta)\tan\theta + 3\ln|M_N/M_E| \right] \quad (4.1)$$

where

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

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

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

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

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

Results for the (4.2) decay mode are presented in Fig. 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 lefthanded neutrinos, mixing between light and heavy neutrinos continues to be suppressed by a factor $s/b \sim O(10^{-6})$, so that the mixing of the light neutrinos among themselves is unitary to a good approximation. Thus the usual Z coupling is essentially flavor conserving and obeys universality. This unitarity also ensures that lefthanded charged-current

radiative contributions to neutral-current processes like $\mu \to e\gamma$ are suppressed by the usual GIM-type mechanism.

For righthanded neutrinos the mixing of the light neutrinos is unitary only up to corrections of order $a/b \sim O(10^{-3,4})$, but the resulting departure from universality of the Z' coupling is insignificant due to the large mass of the Z'. Also the righthanded 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 righthanded 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 lefthanded 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, righthanded 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 inter-actions, flavor-changing and otherwise, are suppressed. The two light Higgs include

actions, flavor-changing and otherwise, are suppressed. The two light Higgs 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 $O(100 GeV/\Lambda_R)$, and again its effects would not have been seen.

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Figure Captions

Figure 1:   Decay rate for $W \to E\overline{N}$ 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.

Figure 2:   Decay rate for Z into two mirror leptons. The solid (dotted) line represents the $N\overline{N}$ ($E\overline{E}$) channel.

Figure 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.

Figure 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.
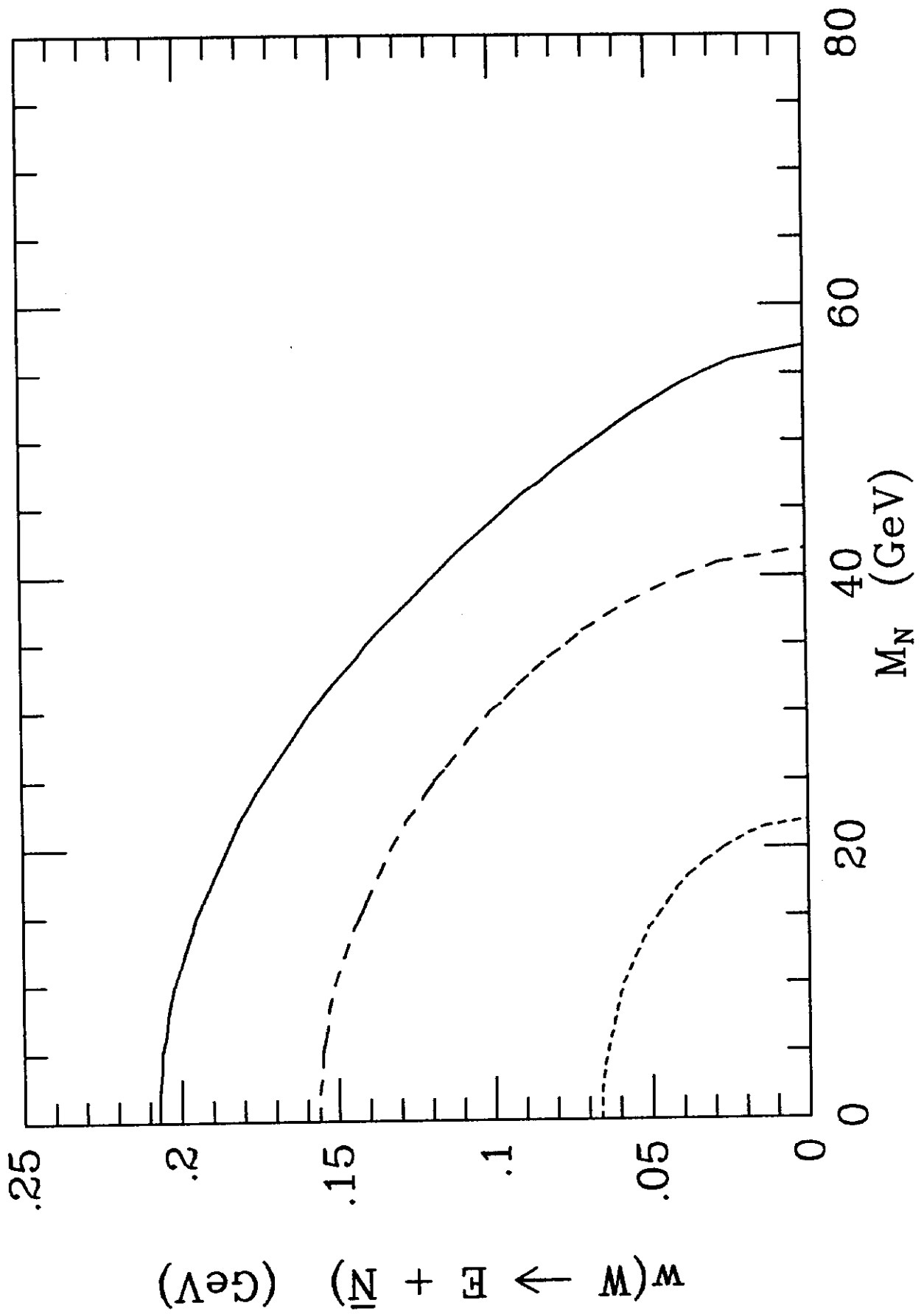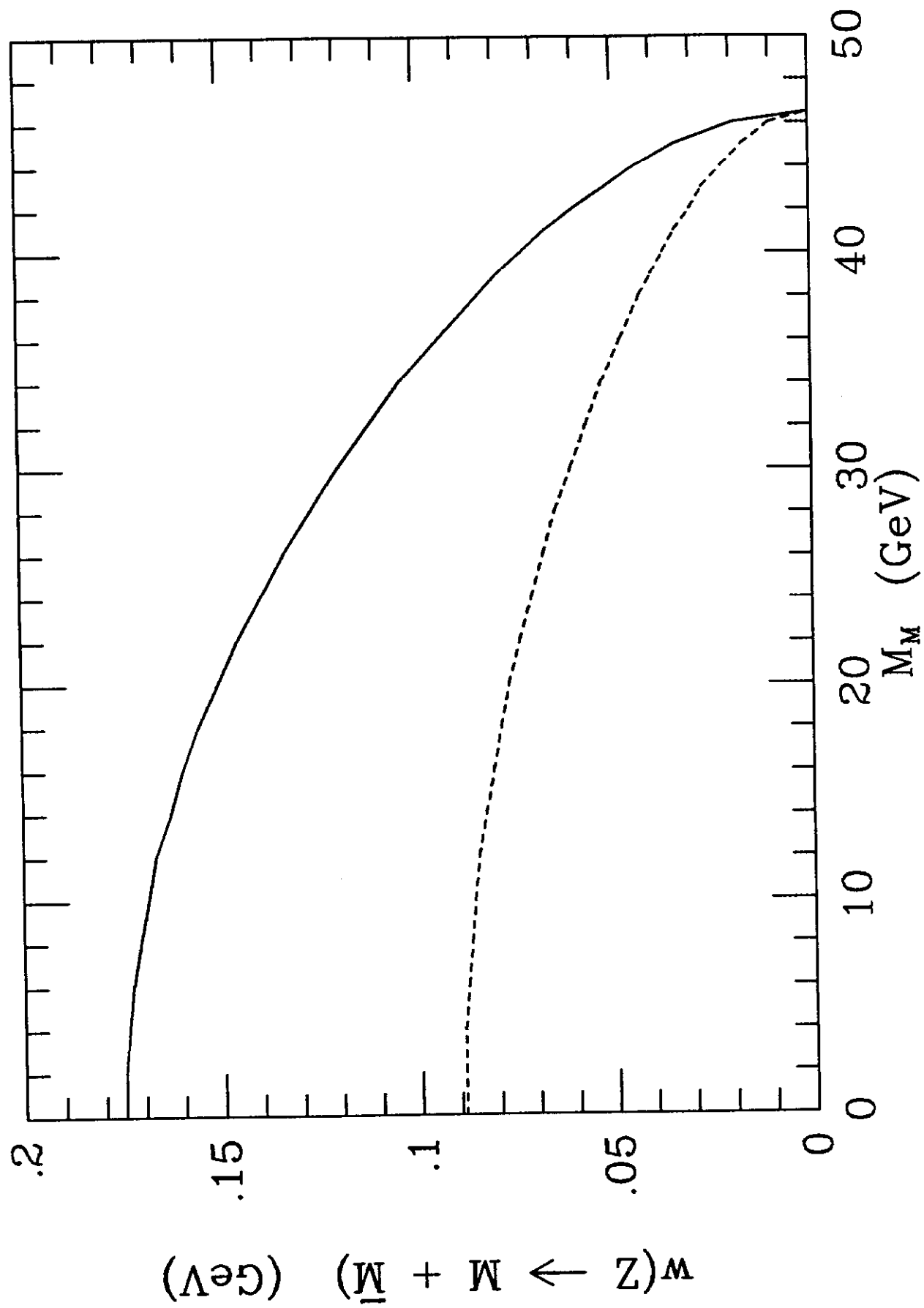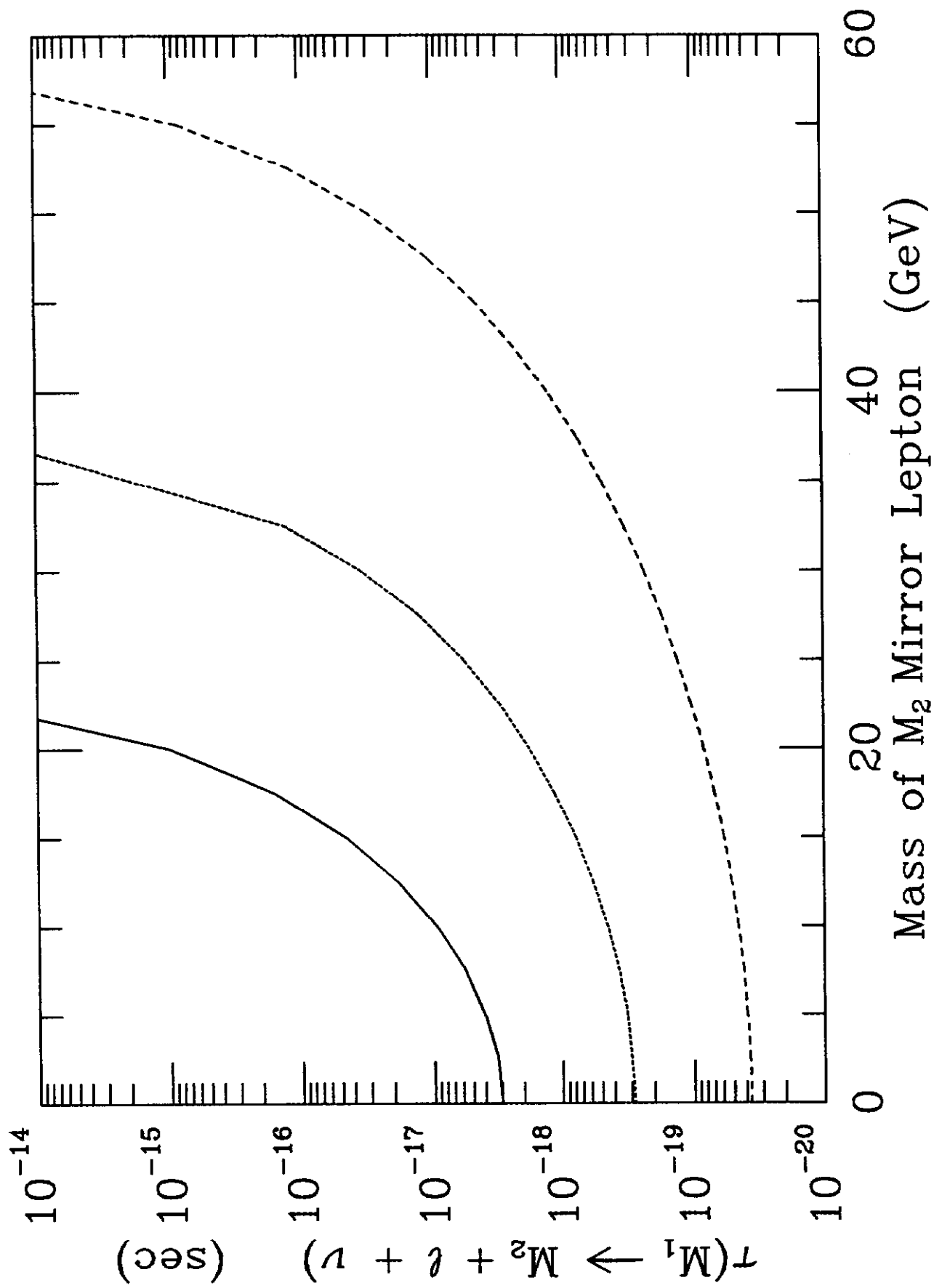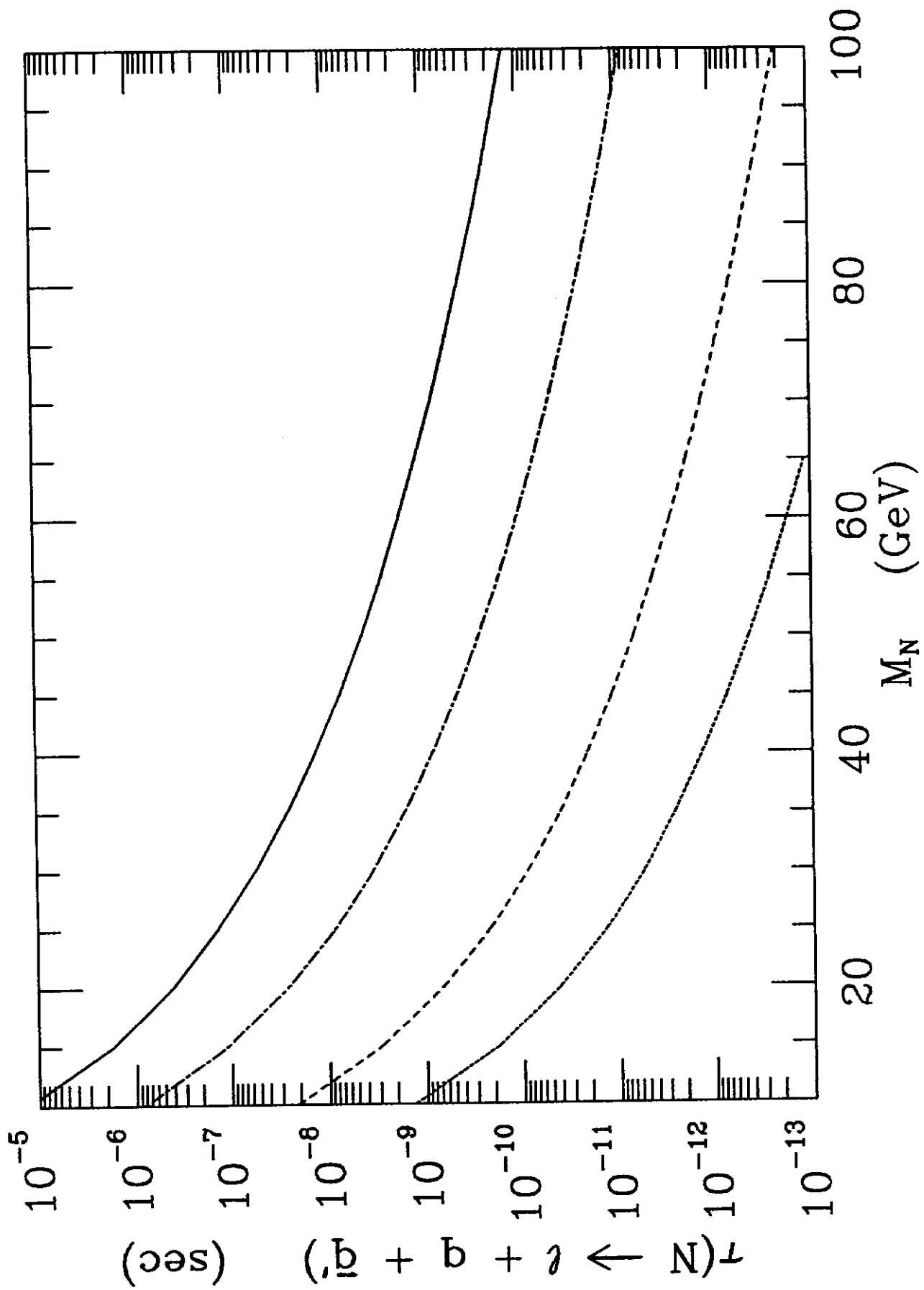
Fig. 1

Fig. 2

Fig. 3

$\tau(M_1 \longrightarrow M_2 + \ell + \nu)$ (sec)

Mass of M_2 Mirror Lepton (GeV)

Fig. 4