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Heavy-Neutrino Chirality Enhancement of the Decay $K_L \to e\mu$ in Left-Right Symmetric Models

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ABSTRACT

We study the decay $K_L \to e\mu$ in minimal extensions of the Standard Model based on the gauge groups $SU(2)_L \otimes U(1)_Y$ and $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$, in which heavy Majorana neutrinos are present. In $SU(2)_L \otimes U(1)_Y$ models with chiral neutral singlets, $B(K_L \to e\mu)$ cannot be much larger than 5×10^{-15} without violating other low-energy constraints. In $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ models, we find that heavy-neutrino-chirality enhancements due to the presence of left-handed and right-handed currents can give rise to a branching ratio close to the present experimental limit $B(K_L \to e\mu) < 3.3 \times 10^{-11}$.

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One of the salient features of the minimal Standard Model (SM) is that the separate leptonic quantum numbers are conserved to all orders of perturbation theory. However, if the SM is considered to be the low-energy limit of a more fundamental theory (*e.g.*, superstrings or grand unified theories), one may then have to worry about large flavour-changing-neutral current (FCNC) effects that could violate experimental data. Among the possible FCNC decays forbidden in the SM, the decay $K_L \to e\mu$ can play a central rôle either to constrain or establish new physics beyond the SM. Furthermore, it is already known that extensions of the SM containing more than one neutral isosinglet can dramatically relax the severe constraints on the mixings between light and heavy neutrinos [1–4], which are dictated by usual see-saw scenarios [5]. Such models with enhanced light-heavy neutrino mixings are also associated with large Dirac mass terms in the general neutrino mass matrix [4]. As an immediate phenomenological consequence, nondecoupling virtual effects originating from heavy neutrinos can considerably enhance the decay rates $H \to l\bar{l'}$ [6], $Z \to l\bar{l'}$ [7], $\tau \to eee$ [8], and the values of other observables [9] to an experimentally accessible level.

In this note, we will investigate similar nondecoupling effects coming from heavy neutrinos and the heavy top quark in the decay $K_L \to e\mu$. We will analyze such effects in Majorana neutrino models based on the gauge groups $SU(2)_L \otimes U(1)_Y$ and $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$ [10–12]. In particular, in general $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$ models with no manifest or pseudomanifest left-right symmetry, a significant enhancement occurs due to the chirality difference between left-handed and right-handed currents, leading to observable rates for the decay $K_L \to e\mu$. The latter may be related to the observation made by Paschos in [13] for the $K_L - K_S$ mass difference in left-right symmetric models.

In order to briefly describe the electroweak sector of the SM with one right-handed neutrino per family, we will adopt the notations given in [4]. Later on, we will extend our analysis to the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge group. To be specific, we will first consider an $SU(2)_L \otimes U(1)_Y$ symmetric model with n_G generations of charged leptons l_i ($i = 1, \dots, n_G$) and light (heavy) Majorana neutrinos ν_α (N_α). The charged-current interaction of this models is then governed by the Lagrangian

$$\mathcal{L}^W_{int} \;=\; -\frac{g_w}{\sqrt{2}} W^{-\mu} \sum_{i=1}^{n_G} \sum_{\alpha=1}^{2n_G} \bar{l}_i \gamma_\mu \mathrm{P}_L B_{l_i\alpha} n_\alpha \quad + \quad H.c., \tag{1}$$

where g_w is the $SU(2)_L$ weak coupling constant, P_L (P_R) $= (1 - (+)\gamma_5)/2$, and B is an $n_G \times 2n_G$ matrix obeying a number of useful identities given in [4]. In Eq. (1), we have collectively defined the mass eigenstates of the light and heavy neutrinos as follows: $n_\alpha \equiv \nu_\alpha$ for $\alpha = 1, \ldots, n_G$ and $n_\alpha \equiv N_{\alpha - n_G}$ for $\alpha = n_G + 1, \ldots, 2n_G$.

In the model under consideration, the decay $K_L \to e\mu$ is induced by the diagrams shown in Fig. 1(a)–(d). In Fig. 1(b)–(d), the field χ_L describes the would-be Goldstone boson in the Feynman–'t Hooft gauge, which is related to the longitudinal polarization of the W boson in the unitary gauge. In the applicable limit of vanishing external momenta, the amplitude of the decay process $K_L \to e\mu$ takes the general form

$$\begin{aligned} A &= \left(\frac{g_w}{\sqrt{2}}\right)^4 \left\langle 0 \left| \bar{d}\gamma_\kappa \mathrm{P}_L s \right| \overline{K}^0 \right\rangle \bar{u}_\mu \gamma^\kappa \mathrm{P}_L v_e \frac{1}{(4\pi)^2} \frac{1}{M_W^2} \\ &\quad \times \sum_{u_i = u,c,t} V_{id}^* V_{is} \sum_{\alpha=1}^{2n_G} B_{\mu\alpha} B_{e\alpha}^* I(\lambda_i, \lambda_\alpha), \end{aligned} \tag{2}$$

where $\lambda_\alpha = m_{n_\alpha}^2 / M_W^2$, $\lambda_i = m_{u_i}^2 / M_W^2$, V is the usual Cabbibo-Kobayashi-Maskawa (CKM) matrix, and the loop function I is obtained from the Feynman graphs in Fig. 1(a)–(d) [16]. The function I is analytically given by

$$I(\lambda_i, \lambda_\alpha) = \left(1 + \frac{1}{4}\lambda_i \lambda_\alpha\right) I_1(\lambda_i, \lambda_\alpha) + 2\lambda_i \lambda_\alpha I_2(\lambda_i, \lambda_\alpha), \tag{3}$$

with

$$I_1(\lambda_i, \lambda_\alpha) = \left[\frac{\lambda_i^2 \ln \lambda_i}{(\lambda_\alpha - \lambda_i)(1 - \lambda_i)^2} + \frac{\lambda_\alpha^2 \ln \lambda_\alpha}{(\lambda_i - \lambda_\alpha)(1 - \lambda_\alpha)^2} - \frac{1}{(1 - \lambda_i)(1 - \lambda_\alpha)}\right], \tag{4}$$

$$I_2(\lambda_i, \lambda_\alpha) = -\left[\frac{\lambda_i \ln \lambda_i}{(\lambda_\alpha - \lambda_i)(1 - \lambda_i)^2} + \frac{\lambda_\alpha \ln \lambda_\alpha}{(\lambda_i - \lambda_\alpha)(1 - \lambda_\alpha)^2} + \frac{1}{(1 - \lambda_i)(1 - \lambda_\alpha)}\right]. \tag{5}$$

Following closely Ref. [14], we define a reduced amplitude $\tilde{A}$ through the expression

$$A = \left(\frac{g_w}{\sqrt{2}}\right)^4 \left\langle 0 \left| \bar{d}\gamma_\kappa \mathrm{P}_L s \right| \overline{K}^0 \right\rangle \bar{u}_\mu \gamma^\kappa \mathrm{P}_L v_e \frac{1}{(4\pi)^2} \frac{1}{M_W^2} \tilde{A}, \tag{6}$$

where

$$\tilde{A} = \sum_{i=u,c,t} V_{id}^* V_{is} \sum_{\alpha=1}^{2n_G} B_{\mu\alpha} B_{e\alpha}^* I(\lambda_i, \lambda_\alpha). \tag{7}$$

The advantage of this definition is that $\tilde{A}$ is a dimensionless quantity carrying the whole electroweak physics. The matrix B obeys a generalized GIM identity [15]

$$\sum_{\alpha=1}^{2n_G} B_{l_1\alpha} B_{l_2\alpha}^* = \delta_{l_1 l_2}, \tag{8}$$

similar to the known one satisfied by the CKM matrix V. A double GIM mechanism is then operative both for the intermediate u-type quarks and neutral leptons, which simplifies Eq. (7) to

$$\tilde{A} = \sum_{i=c,t} V_{id}^* V_{is} \sum_{\alpha=n_G+1}^{2n_G} B_{\mu\alpha} B_{e\alpha}^* E(\lambda_i, \lambda_\alpha), \tag{9}$$

with

$$\begin{aligned} E(\lambda_i, \lambda_\alpha) &= \lambda_i \lambda_\alpha \left\{ -\frac{3}{4} \frac{1}{(1-\lambda_i)(1-\lambda_\alpha)} \right. \\ &+ \left[\frac{1}{4} - \frac{3}{2} \frac{1}{\lambda_i - 1} - \frac{3}{4} \frac{1}{(\lambda_i - 1)^2} \right] \frac{\ln \lambda_i}{\lambda_i - \lambda_\alpha} \\ &+ \left. \left[\frac{1}{4} - \frac{3}{2} \frac{1}{\lambda_\alpha - 1} - \frac{3}{4} \frac{1}{(\lambda_\alpha - 1)^2} \right] \frac{\ln \lambda_\alpha}{\lambda_\alpha - \lambda_i} \right\}. \end{aligned} \tag{10}$$

In Eq. (9), we have considered that the up quark u and the light neutrinos ν_e, ν_μ, and ν_τ are massless. Only the virtual c and t quarks, and the n_G heavy Majorana neutrinos will then contribute to $\tilde{A}$. For definiteness, we have restricted ourselves to a model with $n_G = 2$ (neglecting mixings due to ν_τ), where the mixings B_{lN_α} and the two heavy neutrino masses, m_{N_1} and m_{N_2}, satisfy the relation [7]

$$B_{lN_2} B_{l'N_2}^* = \frac{m_{N_1}}{m_{N_2}} B_{lN_1} B_{l'N_1}^*. \tag{11}$$

The branching ratio for $K_L \to e\mu$ may conveniently be calculated by using isospin invariance relations between the decay amplitudes of $\bar{K}^0 \to \mu^- e^+$ and $K^- \to \mu^- \nu_\alpha$. Setting $m_e = 0$ relative to m_μ in the phase space, one finds

$$B(K_L \to e\mu) = 4.1 \times 10^{-4} \left| \tilde{A} \right|^2. \tag{12}$$

Experimental bounds coming from the nonobservation of the decay $\mu \to e\gamma$ will constrain the parameter space of our theory and impose severe limits on the decay $K_L \to e\mu$. In Majorana neutrino models, the branching ratio of $\mu \to e\gamma$ is

$$B(\mu \to e\gamma) = \frac{6\alpha}{\pi} \Big| \sum_{\alpha=1,2} B_{\mu N_\alpha} B_{eN_\alpha}^* F(\lambda_\alpha) \Big|^2, \tag{13}$$

where the loop function F calculated in [16,17] is given by

$$F(\lambda_\alpha) = \frac{2\lambda_\alpha^3 + 5\lambda_\alpha^2 - \lambda_\alpha}{4(1-\lambda_\alpha)^3} + \frac{3\lambda_\alpha^3 \ln \lambda_\alpha}{2(1-\lambda_\alpha)^4}. \tag{14}$$

The present experimental upper limit $B(\mu \to e\gamma) < 4.9 \times 10^{-11}$ together with Eqs. (11) and (13) can be used to obtain combined constraints on the mixings B_{lN_α} and heavy neutrino masses. These constraints are quite useful in order to individually evaluate the contribution of the charm and top quark to $B(K_L \to e\mu)$. In our numerical analysis, we have used the maximally allowed values $V_{td} = 0.018$ and $V_{ts} = 0.054$, and the central value for the top-quark $m_t = 175$ GeV as has recently been reported by the CDF collaboration [18]. Contrary to [19] where only one heavy neutrino family with mass not much heavier than M_W was considered, we find that the charm-quark contribution is negligible and only top-quark quantum effects are of interest here for heavy neutrino masses larger than 150 GeV. Of course, the mass of the heavy neutrinos should not exceed an upper limit that invalidates perturbative unitarity. This mass limit is qualitatively estimated to be no bigger than 50 TeV [4]. From Fig. 2, we see that the branching ratio takes the maximum value $B(K_L \to e\mu) = 5.5 \times 10^{-15}$, which is still rather far from the present experimental sensitivity $B(K_L \to e\mu) < 3.3 \times 10^{-11}$ at 90% C.L. [20]. In Fig. 2, we have further assumed for the two heavy neutrinos, N_1 and N_2, to have about the same mass m_N. Nevertheless, in Fig. 3, we have plotted the dependence of the branching ratio as a function of the value $\rho = m_{N_2}/m_{N_1}$ for selected values of m_{N_1}. The solid curve in Fig. 3 determines an upper limit of the allowed region for $B(K_L \to e\mu)$ by taking into account the combined constraints arising from the neutrino mixings B_{lN_α} and the validity of perturbative unitarity.

An enhancement of the decay rate can be obtained [14] by considering a general left-right symmetric model based on the gauge group $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$. This model predicts two charged gauge bosons W_L and W_R, which are generally not mass eigenstates, but the relevant mixing angle is proportional to the vacuum expectation value (v_L) of the left-handed Higgs triplet Δ_L with quantum numbers $(0, 1, 2)$. For simplicity, we will work out the realistic case (d) of Ref. [21], in which $v_L = 0$. In this case, the W_L and W_R bosons become mass eigenstates with masses $M_L = M_W$ and M_R, respectively.

In the context of left-right scenarios, there exist three different sets of diagrams depending on the way that the virtual gauge bosons W_L and W_R are involved. To be precise, we group the four diagrams in Fig. 1(a)–(d) separately which are entirely mediated by W_L bosons and are identical with those considered above. As a distinctive set, we consider the Feynman graphs in Fig. 1(e)–(l), in which a W_L and a W_R boson are simultaneously present.

Furthermore, there are four graphs (not shown in Fig. 1) that can be obtained from the first group by replacing the W_L boson by the W_R one. In addition, at tree level, Higgs scalars with FCNC couplings should be taken into account to cancel possible gauge-dependent terms arising from the graphs in Fig. 1 [22]. The inclusion of such Higgs-dependent graphs is not expected to alter quantitatively the results obtained in our analysis. For a discussion on related issues, the reader is referred to Ref. [14,23].

Since the first set of graphs has already been considered in the context of the SM with right-handed neutrinos, we will therefore proceed with the calculation of the graphs of the second one, which are depicted in Fig. 1(e)–(l). Their contribution to the corresponding reduced amplitude is found to be

$$
\begin{aligned}
\tilde{A}_{LR} &= \beta_g\eta \Bigg\{ \sum_{i=c,t} V^{L*}_{id} V^{R}_{is} \sum_{\alpha=n_G+1}^{2n_G} B^{R}_{\mu\alpha} B^{L*}_{e\alpha} (\lambda_i\lambda_\alpha)^{1/2} \\
&\times \left[\left(1+\frac{\beta\lambda_i\lambda_\alpha}{4}\right) J_1(\lambda_i,\lambda_\alpha,\beta) - \frac{1+\beta}{4} J_2(\lambda_i,\lambda_\alpha,\beta)\right] \\
&+ \sum_{i=c,t} V^{R*}_{id} V^{L}_{is} \sum_{\alpha=n_G+1}^{2n_G} B^{L}_{\mu\alpha} B^{R*}_{e\alpha} (\lambda_i\lambda_\alpha)^{1/2} \\
&\left[\left(1+\frac{\beta\lambda_i\lambda_\alpha}{4}\right) J_1(\lambda_i,\lambda_\alpha,\beta) - \frac{1+\beta}{4} J_2(\lambda_i,\lambda_\alpha,\beta)\right] \Bigg\} . \qquad (15)
\end{aligned}
$$

Here, $\beta = M_L^2/M_R^2$, $\beta_g = (g_R^2/g_L^2)M_L^2/M_R^2$, with $g_L = g_w$ and g_R being the coupling constants related to the gauge groups $SU(2)_L$ and $SU(2)_R$, respectively. The parameter η in Eq. (15) is an enhancement factor that results from the different type of operator describing the kaon-to-vacuum matrix element. The factor η is defined as [14]

$$
\begin{aligned}
\eta &= \frac{\left\langle 0 \left| \bar{d}\gamma^\sigma\gamma^\kappa \mathrm{P}_R s \right| \overline{K}^0 \right\rangle \bar{u}_\mu \gamma_\sigma\gamma_\kappa \mathrm{P}_L v_e}{\left\langle 0 \left| \bar{d}\gamma^\alpha \mathrm{P}_L s \right| \overline{K}^0 \right\rangle \bar{u}_\mu \gamma_\alpha \mathrm{P}_L v_e} \\
&\simeq \frac{4M_K^2}{(m_s+m_d)m_\mu} \simeq 50 , \qquad (16)
\end{aligned}
$$

which is estimated by using the assumption of partial conservation of the axial vector current (PCAC). The box functions J_1 and J_2 are given by

$$
\begin{aligned}
J_1(\lambda_i,\lambda_\alpha,\beta) &= \frac{\lambda_i \ln\lambda_i}{(1-\lambda_i)(1-\beta\lambda_i)(\lambda_\alpha-\lambda_i)} + \frac{\lambda_\alpha \ln\lambda_\alpha}{(1-\lambda_\alpha)(1-\beta\lambda_\alpha)(\lambda_i-\lambda_\alpha)} \\
&+ \frac{\beta\ln\beta}{(1-\beta)(1-\beta\lambda_i)(1-\beta\lambda_\alpha)} ,
\end{aligned}
$$

$$J_2(\lambda_i,\lambda_\alpha,\beta) = \frac{\lambda_i^2\ln\lambda_i}{(1-\lambda_i)(1-\beta\lambda_i)(\lambda_\alpha-\lambda_i)} + \frac{\lambda_\alpha^2\ln\lambda_\alpha}{(1-\lambda_\alpha)(1-\beta\lambda_\alpha)(\lambda_i-\lambda_\alpha)} + \frac{\ln\beta}{(1-\beta)(1-\beta\lambda_i)(1-\beta\lambda_\alpha)}\,. \tag{17}$$

In Eq. (15), B^L and V^L are essentially the matrices B and V which have been defined above in the SM with right-handed neutrinos. By analogy, the matrices B^R and V^R parametrize the interaction of W_R with leptons and quarks, respectively. In case of no manifest or pseudomanifest left-right symmetry, there are no experimental constraints on the elements of V^R and they are limited simply by unitarity. In fact, for specific forms of V^R given in Table II of Ref. [14], the $K_L - K_S$ mass difference imposes a lower bound on the W_R-boson mass $M_R \gtrsim 400$ GeV, not very different from the experimental one [20].

Following a similar procedure, we can extract constraints on B^R from the experimental limit of the decay $\mu \to e\gamma$ that can be mediated by right-handed currents. In particular, if the mixings $B^L_{\mu N_\alpha}$ are assumed to be extremely suppressed or vanish, then only W_R bosons can provide a nonzero value to the decay $\mu \to e\gamma$. Therefore, in the evaluation of $B(K_L \to e\mu)$, we will consider only the first term of the r.h.s. of Eq. (15), which is rather conservative. Then, the dominant contribution to the amplitude originates from the diagram (h) in Fig. 1. In the limit of $m_N, M_R \gg M_L$, the reduced amplitude $\tilde{A}_{LR}$ behaves asymptotically as

$$\tilde{A}_{LR} \approx (1.2\ 10^{-4}) \times \eta\beta_g V_{td}\left(\frac{s_L^{\nu_e}}{0.01}\right)\frac{\beta^{1/2}\lambda_t^{3/2}\lambda_N^{1/2}}{4(1-\beta\lambda_t)}\left(\ln\beta + \frac{\lambda_t\ln\lambda_t}{\lambda_t - 1}\right), \tag{18}$$

where constraints for the mixing matrices B^L and B^R coming from the decay $\mu \to e\gamma$ have been implemented. In Eq. (18), we have identified $(s_L^{\nu_l})^2 \equiv \sum_{\alpha=1}^{n_G} |B^L_{lN_\alpha}|^2$, where l stands for an electron or a muon. In agreement with a global analysis of low-energy and LEP data [24], we have considered $(s_L^{\nu_e})^2 = (s_L^{\nu_\mu})^2 \leq 10^{-4}$ in our numerical estimates. Nevertheless, one may have to worry that diagrams similar to Fig. 1(h) and 1(l), which are present in the decay $\mu \to eee$, could lead to a violation of the experimental bound $B(\mu \to eee) < 10^{-12}$ [20]. Considering only the dominant nondecoupling terms, we have estimated that this happens when $\eta m_t V_{td}/(s_L^{\nu_e} m_N) < 1$ for $M_R \sim 1$ TeV. Unless the mass of heavy neutrinos $m_N > 10$ TeV for $M_R \lesssim 1$ TeV, the limits derived from the nonobservation of $\mu \to e\gamma$ will be rather sufficient to preform our combined analysis.

As can be seen from Fig. 4, $B(K_L \to e\mu)$ depends strongly on M_R via the parameter

β. As a natural choice, we have assumed the left-right symmetric case $\beta = \beta_g$. Taking the constraints coming from $\mu \to e\gamma$ into account, we find that heavy neutrinos with few TeV masses can give rise to branching ratios of the order of 10^{-11} close to the present experimental limit. Note that there is a local maximum in Fig. 4 for smaller values of β, where W_R bosons with several TeV masses can also account for $B(K_L \to e\mu) \sim 10^{-11}$.

There is a third set of graphs (not shown in Fig. 1) contained in the reduced amplitude $\tilde{A}_{RR}$, in which a W_L (χ_L) boson should be replaced by a W_R (χ_R) one in Fig. 1(a)–(d). From Eq. (9), it is straightforward to obtain the analytic expression for $\tilde{A}_{RR}$ by making the obvious substitutions mentioned above. In this way, one has

$$\tilde{A}_{RR} = \beta_g^2 \sum_{i=c,t} V_{id}^{R*} V_{is}^{R} \sum_{\alpha=n_G+1}^{2n_G} B_{\mu j}^{R} B_{ej}^{R*} E(\beta\lambda_i, \beta\lambda_\alpha). \tag{19}$$

From Eq. (19), we find numerically that $|\tilde{A}_{RR}|^2$ is about one order of magnitude smaller than $|\tilde{A}_{LR}|^2$ due to the severe constraints coming from the nonobservation of $\mu \to e\gamma$. For instance, for $M_R = 800$ GeV and $m_N = 20$ TeV, the branching ratio $B(K_L \to e\mu)$ is here 1.6×10^{-12} as compared to the value $B(K_L \to e\mu) = 2.0 \times 10^{-11}$ in a complete computation.

In conclusion, we have analyzed some interesting aspects of the decay $K_L \to e\mu$ in the framework of Majorana-neutrino models based on the gauge groups $SU(2)_L \otimes U(1)_Y$ and $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$. In comparison to a previous work [14], we wish to stress that in the renormalizable gauge, diagrams with would-be Goldstone bosons are indeed important, since the stringent limits on the mixings between light and heavy neutrinos can be relaxed by the presence of two heavy neutrino families. In an $SU(2)_L \otimes U(1)_Y$ model with right handed neutrinos, we have found that $B(K_L \to e\mu) \lesssim 10^{-15}$ for heavy neutrinos with TeV masses, where the top-quark contribution prevails over the charm-quark one. In an $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$ model with Majorana neutrinos, the chirality changing diagrams (h) and (l) in Fig. 1 dominate in the branching ratio over the remaining set of graphs. Through heavy-neutrino-chirality enhancements the resulting branching ratio $B(K_L \to e\mu)$ can be as large as the present experimental limit $\sim 10^{-11}$, for a wide range of parameter values. In addition, the constraints derived from $B(K_L \to e\mu)$ are complementary to the ones determined by other low-energy experiments, such as the possible decays $\mu \to e\gamma$ and $\mu \to eee$, and the $K_L - K_S$ mass difference. Therefore, experimental tests at DAΦNE or in other kaon factories will be very crucial and may reveal surprises in the leptonic decay channels of the K_L meson that might signal the onset of new physics.

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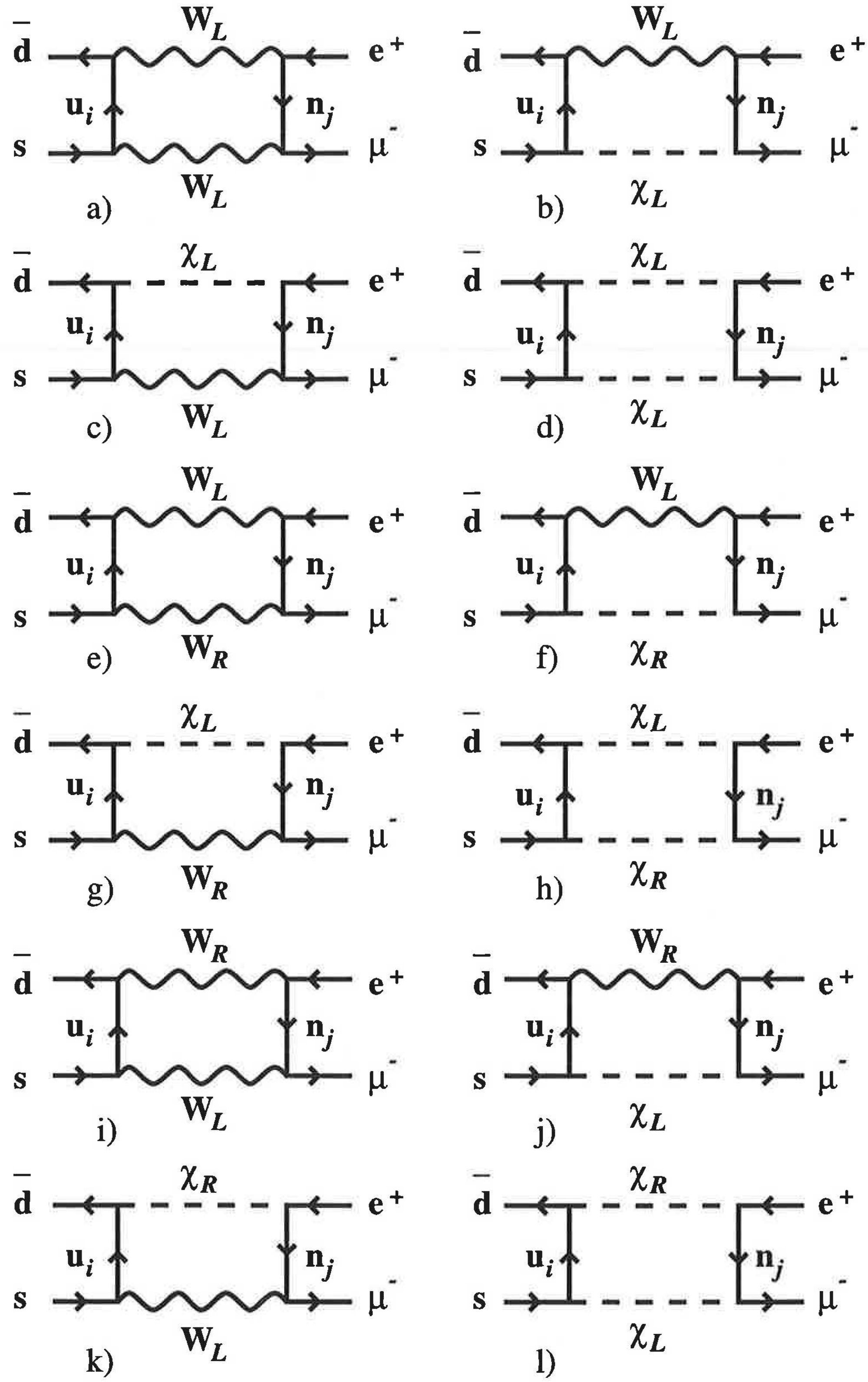
Figure Captions

Fig. 1: Feynman diagrams contributing to $K_L \to e\mu$ in Majorana neutrino models relying on the gauge groups: (a)–(d) $SU(2)_L \otimes U(1)_Y$ and (a)–(l) $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$.

Fig. 2: $B(K_L \to e\mu)$ as a function of the heavy neutrino mass m_N ($\simeq m_{N_1} \simeq m_{N_2}$) ($m_t = 175$ GeV) in the SM with right-handed neutrinos.

Fig. 3: $B(K_L \to e\mu)$ versus $\rho = m_{N_2}/m_{N_1}$ in the $SU(2)_L \otimes U(1)_Y$ model with neutral singlets.

Fig. 4: $B(K_L \to e\mu)$ as a function of $\beta = M_L^2/M_R^2 = \beta_g$ in an $SU(2)_R \otimes SU(2)_L \otimes U(1)_{B-L}$ model, assuming that all heavy neutrinos are approximately degenerate with mass m_N.



d̄
W_L
e+
u_i
n_j
s
μ-
W_L
a)
d̄
W_L
e+
u_i
n_j
s
μ-
χ_L
b)
d̄
χ_L
e+
u_i
n_j
s
μ-
W_L
c)
d̄
χ_L
e+
u_i
n_j
s
μ-
χ_L
d)
d̄
W_L
e+
u_i
n_j
s
μ-
W_R
e)
d̄
W_L
e+
u_i
n_j
s
μ-
χ_R
f)
d̄
χ_L
e+
u_i
n_j
s
μ-
W_R
g)
d̄
χ_L
e+
u_i
n_j
s
μ-
χ_R
h)
d̄
W_R
e+
u_i
n_j
s
μ-
W_L
i)
d̄
W_R
e+
u_i
n_j
s
μ-
χ_L
j)
d̄
χ_R
e+
u_i
n_j
s
μ-
W_L
k)
d̄
χ_R
e+
u_i
n_j
s
μ-
χ_L
l)

Figure 1.

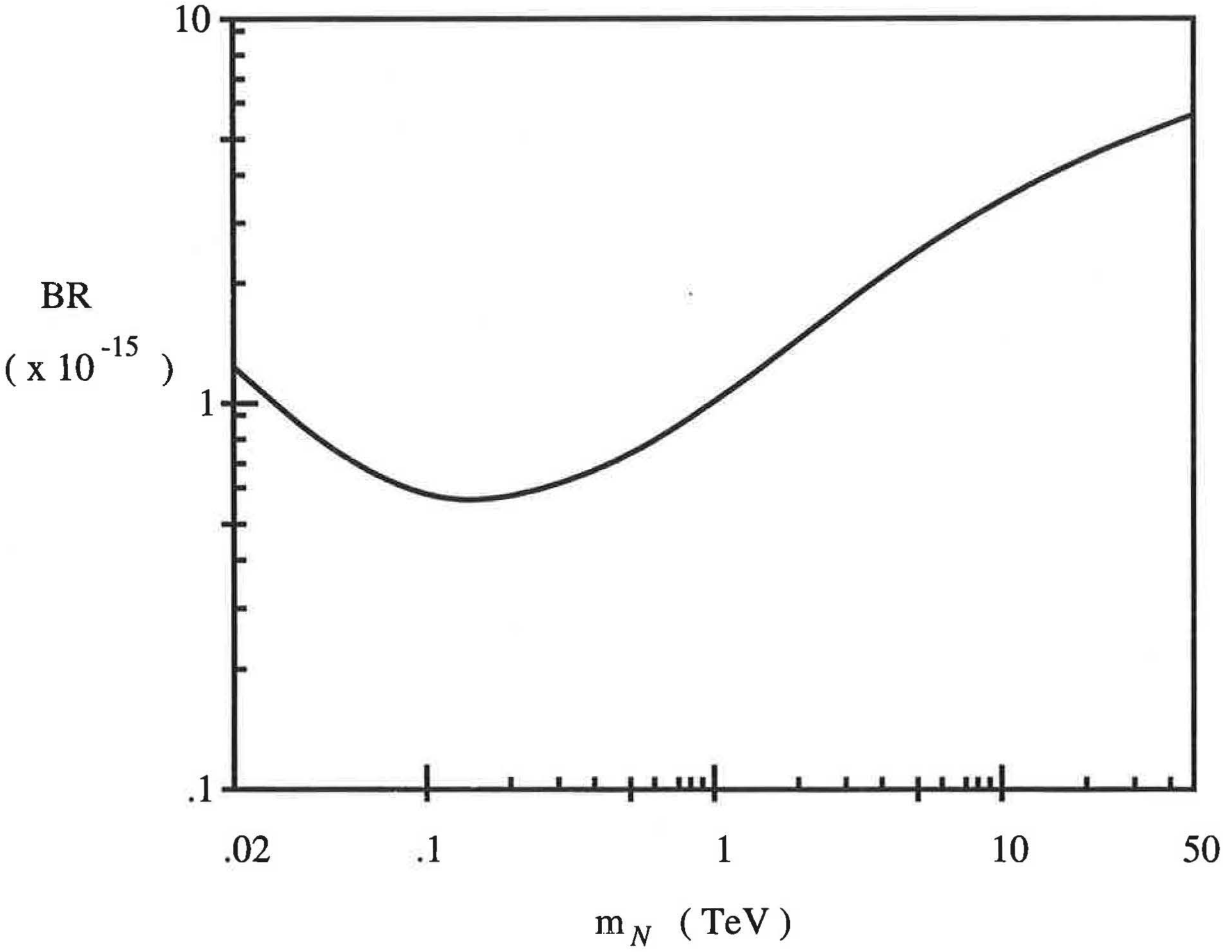
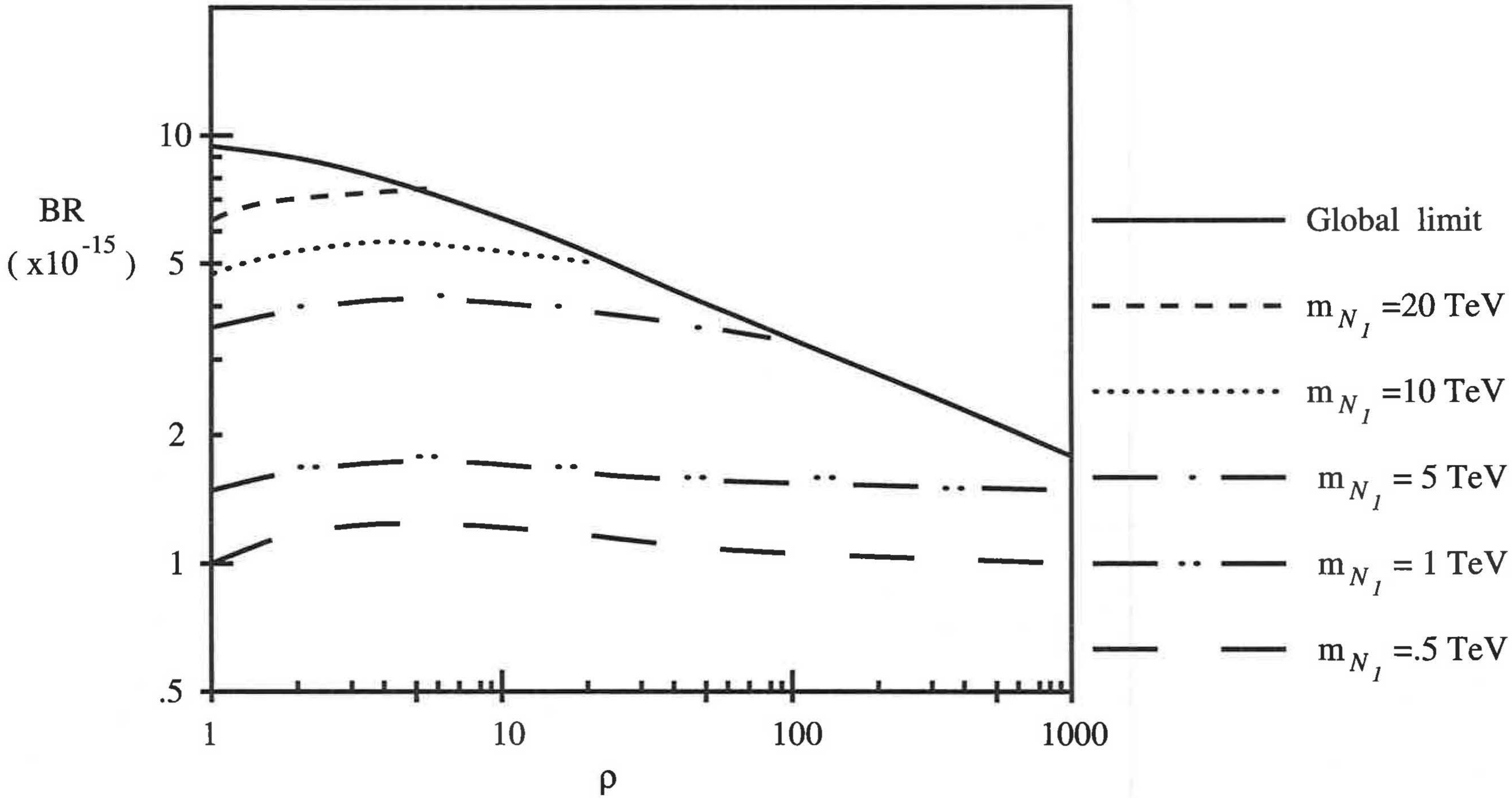


Figure 2.



10
BR
(x10^{-15})
5
2
1
.5
1
10
100
1000
ρ
Global limit
m_{N_1} =20 TeV
m_{N_1} =10 TeV
m_{N_1} = 5 TeV
m_{N_1} = 1 TeV
m_{N_1} =.5 TeV

Figure 3.

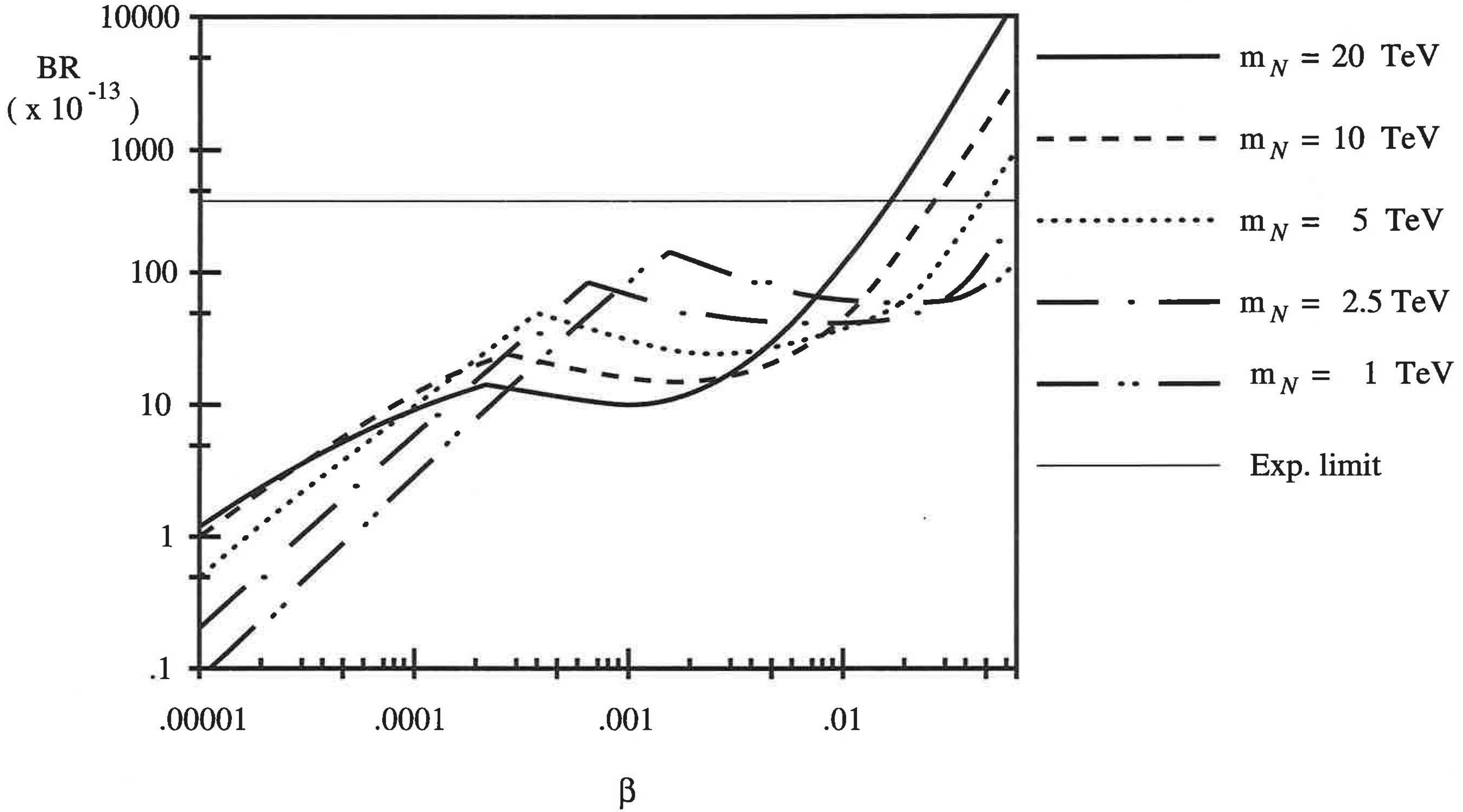

BR
(x 10^-13)
10000
1000
100
10
1
.1
.00001
.0001
.001
.01
β
m_N = 20 TeV
m_N = 10 TeV
m_N = 5 TeV
m_N = 2.5 TeV
m_N = 1 TeV
Exp. limit

Figure 4.