Neutrino mass constraints on \( \mu \)-decay and \( \pi^0 \to \nu\bar{\nu} \)

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In this letter, we show that upper-limits on neutrino mass translate into upper-limits on the class of neutrino-matter interactions that can generate loop corrections to the neutrino mass matrix. We apply our results to \( \mu \)- and \( \pi \)-decays and derive model-independent limits on six of the ten parameters used to parametrize contributions to \( \mu \)-decay that do not belong to the standard model. These upper-limits provide improved constraints on the five Michel parameters, \( \rho, \xi, \xi', \alpha, \alpha' \), that exceed PDG constraints by at least one order of magnitude. For \( \pi^0 \to \nu\bar{\nu} \) we find for the branching ratio: \( B(\pi^0 \to \nu\bar{\nu}) < 10^{-10} \).

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With the discovery of neutrino oscillations a few years ago \( [1, 2, 3] \), the neutrino mass matrix has become a subject of intensive experimental and theoretical research as it provides a unique window into physics beyond the standard model (SM). Indeed, the combination of WMAP \( [4] \), 2DFGRF \( [5] \), and neutrino oscillation data yield an upper limit of 0.23 eV for the mass of an active neutrino. The Planck mission \( [6] \), to be launched in 2007, may further improve this limit to \( \sim 0.04 \) eV \( [7] \). With masses of active neutrinos at least six orders of magnitude smaller than those of all other SM fermions, neutrino masses are presumably generated at an energy scale which significantly exceeds the electroweak scale. At low energies, manifestations of such new physics, including neutrino masses, are suppressed by inverse powers of this heavy scale. For example, in the see-saw mechanism, neutrino masses are inversely proportional to the heavy right-handed neutrino mass, which can range from a few TeVs to \( 10^{15} \) GeV depending on the model.

The study of non-SM neutrino-matter interactions may also shed light on physics beyond the SM. However, since neutrino-matter cross sections are generally small, direct observation of these interactions is experimentally challenging. Moreover, since the number of candidates for physics beyond the SM is large, determining the most viable particle physics scenario is non-trivial. In view of this situation, model-independent constraints on non-SM neutrino-matter interactions in combination with the study of the neutrino mass matrix should prove a valuable tool in the search for new physics.

In this letter we point out a general connection between the neutrino mass and scalar, pseudoscalar, and tensor (S,PS,T) neutrino-matter interactions. In particular, we show that under minimal assumptions these chirality-changing interactions generate contributions to neutrino mass through loop effects. We do not make any assumption about the dynamical origin of the neutrino mass. Instead, we perform a phenomenological analysis and require that such contributions to the mass not exceed the physical neutrino mass. This allows us to place stringent constraints on chirality-changing neutrino-matter couplings. Our general conclusions are then applied to the SM-forbidden decay of \( \pi^0 \) into a neutrino and an anti-neutrino with the same helicity (\( \pi^0 \to \nu\bar{\nu} \)) and to \( \mu \)-decay. In the former case we show that the cosmological neutrino mass upper-limit constrains the branching ratio for \( \pi^0 \to \nu\bar{\nu} \) to be \( \sim 10^4 \) times smaller than the best current experimental limit \( \pi^0 \to \nu\bar{\nu} \).

General Argument. The general chirality-changing effective neutrino-fermion interaction can be written as

\[
\mathcal{L} = G_F \sum_{l,l',f,f',i} a_{i,l,l',f,f'}^l \bar{\nu}_{i} \Gamma_{l, f} \nu_{l'} + \text{h.c.},
\]

where \( i = S, P, \), with \( \Gamma_{S} = 1 \), \( \Gamma_{P} = \gamma^5 \), and \( \Gamma_{T} = \sigma^{\mu\nu} \); the sum over \( l, l' \) runs over active neutrino flavors while the sum over \( f, f' \) is over the SM charged fermions (this approach does not yield competitive constraints on neutrino-neutrino scattering), and \( a_{i,l,l',f,f'}^l \) are dimensionless constants parametrized in terms of the Fermi constant \( G_F = 1.16637(1) \times 10^{-5} \) GeV\(^{-2} \). The chirality-changing interaction in Eq. \( 1 \) generally contributes to the neutrino mass via diagrams shown in Fig. \( 1 \). Substituting \( \bar{\nu}_i \Gamma_{l, f} \nu_{l'} \) in Eq. \( 1 \) induces Majorana neutrino masses.

Eq. \( 1 \) is a general effective Lagrangian for neutrino-matter interactions constructed from non-renormalizable operators (the coupling constants have negative mass dimension as seen from the overall factor of \( G_F \)). Therefore, a new counterterm will be needed for each operator to cancel divergences that may appear in the eval-
FIG. 1: One- and two-loop contributions to the neutrino mass (denoted \( \delta m^{(1)}_\nu \), and \( \delta m^{(2)}_\nu \) respectively) generated by chirality-changing neutrino-fermion interaction. For Majorana mass (denoted \( \delta m^{(0)}_\nu \)), the mass of fermion \( f \) couple to the weak boson that is inserted, \( \delta m^{(1)}_\nu \) \( \approx \nu (2) \) / \( \delta m^{(2)}_\nu \), where the superscript \( \delta m^{(1)}_\nu \) is order ten as discussed below. Furthermore, the mass dependence of each loop diagram must be an expansion series in powers of \( (m^2_\nu/M^2_Z)^n m_f \) with \( n = 0, 1, 2, \ldots \), and where the \( n = 0 \) term appears only at second order with the exchange of a weak boson as in the diagrams of Fig. 1. It follows that the \( L = 2 \) diagrams with a weak boson are largest except for the case where \( m_f = m_{top} \) as mentioned above. We thus have the counterintuitive result that the 1-loop graph is generally subdominant.

The \( (\ln \mu^2)^2 \) factors in Eq. (2) appear because the diagrams are ultraviolet divergent; they are compensated by the \( \mu \)-dependence of the neutrino mass counterterm and the \( \mu \)-dependence of the \( a^{ff'}_{\nu} \)'s deduced from the renormalization group (RG) equations they satisfy. Thus, in order to extract constraints on the \( a^{ff'}_{\nu} \)'s, one must choose a renomalization scale \( \mu \).

This value of \( \mu \) should exceed the mass of the heaviest particle included in the effective field theory (EFT)—in our case \( m_t \), the top quark mass—while at the same time take into account the scale at which the onset of new physics might be expected. We choose the renormalization scale to be around 1 TeV, a scale often associated with physics beyond the SM in many particle physics models. Since \( \mu \) appears in a logarithm, our conclusions do not depend strongly on its precise value. Note that the renormalization scale \( \mu \) is far above the energy scale at which processes like \( \mu \)-decay and \( \pi \rightarrow \nu \nu \) occur. In principle, the couplings appearing in Eq. (2) should be evolved down to \( \sim 1 \) GeV using the appropriate RG equations, but this can at most generate factors of \( O(1) \). For example, the running of coupling constant associated with the four-quark operators in kaon decay, from the weak scale down to \( \mu \approx 1 \) GeV, generates only factors of \( 2 \). There is no reason to expect a more substantial change to the four-lepton or quark-lepton operators of Eq. (1) when running \( \mu \) down to \( \sim 100 \) MeV. Thus, in the model-independent analysis of this letter, there is no need to take the RG running of coupling constants into account. We emphasize that values of \( \mu \) below the weak scale can not be substituted in Eq. (2). Below the weak scale, the dependence of the amplitude on \( \mu \) becomes suppressed by inverse powers of the weak scale as required by the decoupling theorem. See the section on QCD renormalization in Ref. [10] for a more detailed discussion of this point in the case of QCD. Note that \( \mu \) would not appear in a specific model where neutrino masses are calculated radiatively from finite diagrams. In that case, the logarithms would instead have arguments of the form \( M^2_\Lambda/M^2_Z \).}

\[
\delta m^{(1)}_\nu \approx N_c G_F a^{ff'}_{\nu \chi \eta \chi'} m^3_f (\ln \mu^2/m^2_f)/4 \pi^2 \\
\delta m^{(2)}_\nu \approx g^2 N_c G_F a^{ff'}_{\nu \chi \eta \chi'} (m_f or m_f') M^2_Z (\ln \mu^2/M^2_Z)^2 (4\pi^2)^L
\]

where \( L \geq 2 \) is the loop order and the logarithm is of order ten as discussed below. Furthermore, the mass dependence of each loop diagram must be an expansion series in powers of \( (m^2_\nu/M^2_Z)^n m_f \) with \( n = 0, 1, 2, \ldots \), and where the \( n = 0 \) term appears only at second order with the exchange of a weak boson as in the diagrams of Fig. 1. It follows that the \( L = 2 \) diagrams with a weak boson are largest except for the case where \( m_f = m_{top} \) as mentioned above. We thus have the counterintuitive result that the 1-loop graph is generally subdominant.

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Below we use Eq. (2) to constrain \( a^{ff'}_{\nu} \) by requiring \( \delta m_\nu \equiv \delta m^{(1)}_\nu + \delta m^{(2)}_\nu \lesssim m_\nu \) where \( m_\nu \) is the physical neutrino mass. Since the graphs of Fig. 1 are divergent, there will be counterterms that absorb the infinities. In
the absence of fine tuning and assuming perturbation theory to be valid, the leading log contributions of the loop graphs should be no larger than the physical value of neutrino masses.

We now apply our general results to non-SM \( \pi^0 \) and \( \mu \) decays. We adopt the upper limit of 0.7 eV on the sum of the neutrino masses from Ref. [4], which translates into the limit \( m_\nu < 0.23 \) eV for individual neutrino masses when neutrino oscillation constraints are included.

\( \pi^0 \)-decay. We obtain from Eq. (2) \( a_{PS,LL}^{\pi_0} < 10^{-3} \) for \( q = u, d \). For the calculation we used \( m_\pi = m_\mu = m_\tau = (m_u + m_d)/2 = 4 \) MeV [10] (constituent quark masses are inappropriate when working at \( \mu \sim 1 \) TeV).

We can use this result to place an upper limit on the branching ratio \( B(\pi^0 \to \nu \bar{\nu}) \). Starting from the neutrino-quark interaction Lagrangian in Eq. (1) with \( \Gamma_{PS} = \gamma^5 \) we obtain the effective interaction

\[
\mathcal{L}_{\pi^0 \nu \bar{\nu}} = \frac{G_F}{\sqrt{2}} F_\pi m_\pi^2 \left( a_{PS,LL}^{\text{uu}} - a_{PS,LL}^{\text{dd}} \right) \pi^0 \bar{\nu}_i \gamma^\mu \nu_j , \tag{3}
\]

where \( F_\pi = 92.4 \) MeV is the pion decay constant and \( m_\pi \) is the pion mass. The above equation leads to the branching ratio \( B(\pi^0 \to \nu \bar{\nu}) = 10^{-10} \left( a_{PS,LL}^{\text{uu}} - a_{PS,LL}^{\text{dd}} \right)^2 < 10^{-10} \), which is four orders of magnitude stronger than the current best experimental limit \( B(\pi^0 \to \nu \bar{\nu})_{\text{Exp}} < 8.3 \times 10^{-7} \) where \( l = t' = \mu \). Our limit on \( B(\pi^0 \to \nu \bar{\nu}) \) improves by a further two orders of magnitude if the possible Planck limit of \( m_\nu < 0.04 \) eV is used instead of \( m_\nu < 0.23 \) eV.

\( \mu \)-decay. Muon decay can be described with the following effective interaction [10]

\[
\mathcal{L}_{\mu \to e \nu, \bar{\nu}} = \frac{4G_F}{\sqrt{2}} \sum_{\gamma = S, V, T} \sum_{\epsilon = \mu, R, L} \sum_{\xi = 0, 1, 2} g_{\mu \gamma}^{\epsilon \xi} \epsilon_\gamma \nu^\mu \nu^\epsilon \mu_{\text{muon}} \Gamma_{\gamma \mu} , \tag{4}
\]

where \( \gamma = S, V, T \) indicate, respectively, scalar, vector, and tensor interactions and \( \epsilon, \mu = R, L \) indicate the chiralities of the charged leptons. The chiralities \( n \) and \( m \) of the neutrinos are determined by the values of \( \gamma, \epsilon \), and \( \mu \). The constants \( g_{\mu \gamma}^{\epsilon \xi} \) parameterize the strength of the corresponding phenomenological interactions and can be related to the \( a_{\mu \nu}^{\text{LL}} \) through Fierz transformations. In the SM, \( g_{\mu \gamma}^{\epsilon \xi} = 1 \) with the rest being zero.

Limits on \( g_{RL}^{S, V, T} \), \( g_{RL}^{S, V, T} \), \( g_{RL}^{S, V, T} \), \( g_{RL}^{S, V, T} \), and \( g_{RL}^{S, V, T} \) can be obtained from Fig. 1(d) and Eq. (2) with \( M_2 \to M_3 \) and \( M_2 = m_\tau \), the mass of the electron, for \( g_{RL}^{S, V, T} \) and \( g_{RL}^{S, V, T} \). The second column displays current upper-limits from Ref. [10]. Except for \( g_{RL}^{S, V, T} \), our model-independent upper-limits are at least one order of magnitude better than the ones appearing in Ref. [10].

The limits on \( g_{\mu \gamma}^{\epsilon \xi} \) translate into order of magnitude upper-limits on the MPs. Using the definitions in Ref. [10], and their limit on \( (b + b')/A < 10^{-3} \) at 90% CL as well as the fact that \( A \cong 16 \), we obtain the limits given in Tab. II. The meaning of the numbers is explained in the caption. The bracketed limits on \( \xi' \) are not fully constrained by upper-limits on neutrino mass. They are included in the table because the parameters with the largest uncertainties that enter its definition are here better constrained. In particular, the largest uncertainty in

\[
1 - \xi' = [(a + a') + 4(b + b') + 6(c + c')]/A , \tag{5}
\]

stems from the relatively large PDG upper-limits on the parameters \( a, a', c, c' \) when compared to the upper-limit on \( (b + b')/A \). Our limits on the former parameters is substantially better, thus improving on the PDG limit.

| MP          | PDG | WMAP/GRS | TWIST/PSI | Planck/GRS |
|-------------|-----|----------|-----------|------------|
| \( \rho - 3/4 \) | 7   | 1        | 0.1       | 0.1        |
| \( \eta \)   | 33  | X        | 0.1       | X          |
| \( \delta - 3/4 \) | 7   | 10       | 0.1       | 0.1        |
| \( 1 - \xi' \) | 80  | [10]     | X         | [4]        |
| \( 1 - \xi' \) | 58  | 10       | X         | 0.1        |
| \( a/A \)    | 9   | 0.001    | X         | 0.0001     |
| \( a'/A \)   | 8.8 | 0.001    | X         | 0.0001     |
| \( c/A \)    | 15.9 | 1       | X         | 0.1        |
| \( c'/A \)   | 13  | 1        | X         | 0.1        |
| \( c'/A \)   | 7.5 | 0.1      | X         | 0.01       |
| \( c'/A \)   | 7.5 | 0.1      | X         | 0.01       |

TABLE I: Approximate upper-limits on the \( g_{\mu \gamma}^{\epsilon \xi} \)’s from Ref. [10] in comparison to the ones derived from the loop-graph of Fig. 1(d) in combination with cosmological limits on the neutrino mass.

| MP          | PDG | WMAP/GRS | TWIST/PSI | Planck/GRS |
|-------------|-----|----------|-----------|------------|
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| \( \eta \)   | 33  | X        | 0.1       | X          |
| \( \delta - 3/4 \) | 7   | 10       | 0.1       | 0.1        |
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| \( 1 - \xi' \) | 58  | 10       | X         | 0.1        |
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TABLE II: Order of magnitude upper-limits on the MPs. All numbers should be multiplied by \( 10^{-3} \). Note that \( a, a', c, c' \) are not technically MPs, and instead belong to a set of parameters defined by Kinoshita and Sirlin [11]. PDG numbers are given in the second column at 95% confidence level (CL) and 90% CL (numbers with daggers). The third column shows the order of magnitude limits extracted from the \( g_{\mu \gamma}^{\epsilon \xi} \)’s given in Tab. I and Ref. [10]. The fourth column gives expected order of magnitude limits from the TWIST and PSI experiments [11, 12]. The fifth column refers to improved limits on the MPs due to the anticipated data from the Planck Mission expected to constrain the upper-limit on the neutrino mass to around \( m_\nu \lesssim 0.04 \) eV; see also Refs. [11, 17]. The meaning of the bracketed numbers is explained in the text.
does not constrain \((b + b')/A\). With the improved limits on \(a, a', c, c'\) due to Planck data, the upper-limit on \(1 - \xi'\) should then be entirely due to the upper-limit on \(4(b + b')/A\). In the same vein, note that because our constraint on \(\alpha\) is so strong, the measurement of \(\eta = (\alpha - 2\beta)/A\), at PSI [15] to a few parts in \(10^{-4}\) will also constitute a measurement of the MP \(\beta\).

Finally, note that a similar analysis for the decay \(\tau \rightarrow \mu \bar{\nu}_\mu \nu_\tau\) can be performed

\[
\mathcal{L}_{\tau \rightarrow \mu \bar{\nu}_\mu \nu_\tau} = \frac{4G_F}{\sqrt{2}} \sum_{\gamma = \beta, \gamma' = \beta', \mu, \tau = \mu, \tau} g^{(\gamma)}_{\mu, \tau} \tilde{\nu}_\mu \gamma^\nu \bar{\nu}_\tau \gamma \tau\, \nu_\tau, \tag{6}
\]

the following limits are obtained: \(g^{(\beta')}_{\mu, \tau} \tilde{V}(\tau) V^{(\gamma)}(\tau) T(\tau) < 10^{-4}\) and \(g^{(\beta)}_{\mu, \tau} \tilde{V}(\tau) V^{(\gamma)}(\tau) T(\tau) < 10^{-6}\). In a particle physics model where the charged-lepton decay couplings are all of the same order, the \(g^{(\tau)}_{\mu, \tau}\) should provide the best limits on the MPs.

**Non-SM contributions to neutral currents.** In light of the NuTeV result on the weak mixing angle \((\theta_W) [12]\), constraining non-SM neutral currents is particularly timely. To determine \(\sin^2 \theta_W\), the experiment measures the ratio of neutral to charged currents in \(\nu_\mu (\bar{\nu}_\mu)\)-quark interactions. Any deviation from the SM neutral or charged current can be interpreted as a deviation from the SM predictions for \(\sin^2 \theta_W\). For neutral currents, the relevant coupling constants are \(a^{ee}_{\mu, \tau} < 10^{-3}\) for \(q = u, d\) and \(i = S, P, S, T\). The (axial-)vector currents of the SM do not change the chirality or the flavor of the neutrino while the chirality-changing coupling interactions consid- ered in this work do. Thus, the final states are different and the rates—not the amplitudes—must be added. Therefore, the chirality-changing non-SM operators can at most modify the SM neutral current by \(10^{-6}\) and can not account for the NuTeV anomaly.

**Conclusions.** Derivation of our results requires only minimal assumptions. We view the SM as an EFT valid below a certain energy scale (taken to be above \(1\) TeV) and assume the validity of perturbation theory. Note that the interactions of Eq. [11] are not gauge-invariant under \(SU(2)_L \times U(1)_Y\). From a strictly formal point of view, our EFT is not allowed since \(\mu\) is above the weak scale; the operators could be embedded in a gauge-invariant structure, but the resulting Ward identities may impose relationships between the parameters which are assumed independent in this letter. However, since the neutrino mass does not violate gauge-invariance (e.g., in the SM, the neutrino mass is generated through the spontaneous breaking of a gauge symmetry), diagrams that contribute to \(m_e\) are not forced to cancel in a gauge-invariant model. Our order-of-magnitude estimates should therefore be robust—finely tuned cancellations not withstanding. The MPs \(\delta\) and \(\rho\) will soon be constrained with improved precision by the TWIST experiment at the \(10^{-4}\) level [11]. Although results of such measurements will be valuable whether or not a positive signal is observed, an especially interesting situation would arise in the case where TWIST measured finite deviations from the SM values of \(\delta\) and \(\rho\) since that would have implications for the neutrino mass. Thus, any particle physics model that could accommodate deviations of \(\delta\) and \(\rho\) at the \(10^{-3} - 10^{-4}\) level would also be challenged to simultaneously generate neutrino masses consistent with observations; for example, this could be achieved through finely tuned cancellations of the radiative corrections to the neutrino mass shown in Fig. [1] or mixing with right-handed neutrino states with masses \(\approx 0.23\) eV that could lead to large contributions to the MP’s. Furthermore, such a measurement would have implications for all physical processes where the magnitude of the neutrino mass plays a role, like \((\nu_\beta \beta)\)-decay when the neutrino is a Majorana fermion.

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