

RECEIVED: May 7, 2020

REVISED: June 20, 2020

ACCEPTED: July 1, 2020

PUBLISHED: July 28, 2020

Scalar and tensor neutrino interactions

Tao Han,^a Jiajun Liao,^b Hongkai Liu^a and Danny Marfatia^c

^a*Department of Physics and Astronomy, University of Pittsburgh,
3941 O'Hara Street, Pittsburgh, PA 15260, U.S.A.*

^b*School of Physics, Sun Yat-sen University,
135 Xingang Xi Road, Guangzhou, 510275, China*

^c*Department of Physics and Astronomy, University of Hawaii at Manoa,
2505 Correa Road, Honolulu, HI 96822, U.S.A.*

E-mail: than@pitt.edu, liaojiajun@mail.sysu.edu.cn, hol42@pitt.edu,
dmarf8@hawaii.edu

ABSTRACT: We constrain general Dirac neutrino interactions based on the Standard Model Effective Field Theory framework extended with right-handed neutrinos N (SMNEFT) using deep inelastic and coherent elastic neutrino scattering, nuclear beta decay, and meson decay data, and high energy electron-proton and proton-proton collider data. We compute the one-loop anomalous dimensions of the low-energy effective field theory (LEFT) below the electroweak scale and of SMNEFT above the electroweak scale. The tree-level matching between LEFT and SMNEFT is performed at the electroweak scale. Currently, the most stringent limits on scalar and tensor interactions arise from pseudoscalar meson decays and the LHC measurements at the per mille level. In the future, the upcoming High-Luminosity LHC (HL-LHC) has the potential to reach the 10^{-4} level and LHeC can play an important role under certain theoretical assumptions.

KEYWORDS: Neutrino Physics, Effective Field Theories, Beyond Standard Model

ARXIV EPRINT: [2004.13869](https://arxiv.org/abs/2004.13869)

Contents

1	Introduction	1
2	Theoretical formalism	3
2.1	LEFT	3
2.2	SMNEFT	4
2.3	RG running and matching	5
3	Low-energy constraints	7
3.1	Neutrino mass bounds	7
3.2	Pseudoscalar meson decay	8
3.3	Nuclear beta decay	11
3.4	Neutrino deep inelastic scattering	12
3.4.1	CHARM: $\nu_e q \rightarrow \nu q$	13
3.4.2	NuTeV: $\nu_\mu q \rightarrow \nu q$	14
3.5	CE ν NS	15
4	Collider constraints	19
4.1	Proton-proton colliders	20
4.2	Electron-proton colliders	22
5	Summary	23
A	1-loop RG running	25

1 Introduction

Flavor oscillations between massive neutrinos is a firmly established phenomenon that cannot be accounted for by the Standard Model (SM) [1–3], thus strongly motivating physics beyond the SM (BSM) associated with the neutrino sector [4, 5]. The absence of BSM signals at the LHC and other low energy experiments suggests that either the new physics scale is beyond the kinematic reach of our high energy experiments or that it is of a more subtle incarnation. It is therefore prudent to guide our search for new physics as generally as possible without theoretical prejudice.

Generically, departures from the SM at energies below a new physics scale can be described by a model-independent Effective Field Theory (EFT) after integrating out the heavy degrees of freedom in the new physics sector. Such an effective Lagrangian was first constructed by Fermi for the 4-fermion contact interaction involving a neutrino in nuclear β -decay. In the SM language, the contact interaction is a result of integrating

out a heavy particle, the electroweak W -boson. To account for nonstandard interactions (NSI) of neutrinos, Wolfenstein proposed 4-fermion interactions with general couplings [6], that have helped understand matter effects in solar neutrino oscillation experiments. Much theoretical and experimental effort has been made to search for potential new physics along the lines of neutrino NSI; for reviews see refs. [7–9]. To explore new physics near or above the electroweak scale, it is appropriate to adopt an EFT, respecting the full SM gauge symmetry with the SM field content, the so-called Standard Model Effective Field Theory (SMEFT) [10–12]. Neutrino NSIs at low energies can be obtained in this framework by integrating out the heavy particles t, W^\pm, Z and the Higgs boson h .

The addition of right-handed neutrino states to the SM is the simplest extension to account for neutrino oscillations. In this article, we adopt this framework including three right-handed neutrino states N that are lighter than a keV. Naturally, they are “sterile neutrinos” with no SM gauge charges. However, we do not specify their possible Majorana mass terms. We restrict our study to the case in which the left-handed neutrino states are Dirac in nature. Then, new flavor-conserving neutral current vector and tensor interactions are possible. The general Standard Model Effective Field Theory extended with right-handed neutrinos (SMNEFT) has been presented in refs. [13, 14]. We follow this well motivated formalism. A framework for model-independent General Neutrino Interactions (GNI) below the electroweak scale has been also constructed in ref. [15]. All operators of scalar, pseudoscalar, vector, axial vector and tensor interactions of neutrinos with SM fermions are included, leading to potentially rich phenomenology. We note that scalar and tensor GNI operators cannot be embedded in SMEFT at the dimension-six level, but are present in SMNEFT.

There are well-motivated ultraviolet complete models that introduce SM singlet right-handed neutrinos. $U(1)_{B-L}$ extensions of the SM generate vector interactions between right-handed neutrinos and SM particles. In models with left-right symmetry [16–18], such interactions are generated after left-right symmetry breaking. Leptoquark models [14, 19] generate various neutrino-quark interactions. Each of these leads to model-dependent phenomenology if the new states are kinematically accessible, which we will not explore further. Instead, we focus on generic features of Dirac neutrino interactions after integrating out the heavy degrees of freedom.

In this article, we set out to examine new scalar and tensor neutrino-quark interactions using SMNEFT. We present constraints on the operators from current low-energy data including neutrino mass bounds, meson and nuclear decays, neutrino deep inelastic scattering, and coherent elastic neutrino-nucleus scattering ($\text{CE}\nu\text{NS}$), and from current high-energy data from the LHC. We also explore the potential sensitivities for future $\text{CE}\nu\text{NS}$, LHeC, and LHC experiments. Since the physical processes we consider span a vast range of energies, the renormalization group running and matching effects for the relevant Wilson coefficients at different energy scales have been properly taken into account in our analyses; this procedure was carried out for vector NSI in ref. [20].

The rest of the paper is organized as follows. In section 2 we provide an overview of the theoretical formalism and emphasize the importance of matching the Wilson coefficients by renormalization group running at different scales. In section 3, we present the current

constraints and future sensitivities from a wealth of low energy experiments. In section 4, we study the current and projected bounds from high-energy proton-proton colliders and electron-proton colliders. We summarize our results in section 5. Some details of the renormalization group running are given in appendix A.

2 Theoretical formalism

The 4-fermion neutrino interactions via the SM neutral and charged currents at the leading order (LO), after integrating out the Z and W^\pm propagators, are

$$\mathcal{L}_{\text{SM}}^{\text{NC}} = -\frac{G_F}{\sqrt{2}}\delta^{\alpha\beta}\delta^{\gamma\delta}[\bar{\nu}_\alpha\gamma_\mu(1-\gamma_5)\nu_\beta][g_{L,f}\bar{f}_\gamma\gamma^\mu(1-\gamma_5)f_\delta + g_{R,f}\bar{f}_\gamma\gamma^\mu(1+\gamma_5)f_\delta], \quad (2.1)$$

$$\mathcal{L}_{\text{SM}}^{\text{CC}} = -\frac{G_F}{\sqrt{2}}V_{\delta\gamma}^*\delta^{\alpha\beta}[\bar{\nu}_\alpha\gamma_\mu(1-\gamma_5)\ell_\beta][\bar{d}_\gamma\gamma^\mu(1-\gamma_5)u_\delta] + \text{h.c.}, \quad (2.2)$$

where the Fermi constant $G_F/\sqrt{2} = (2v^2)^{-1} = g^2/8M_W^2$. f denotes quarks and charged leptons, V is the CKM quark-mixing matrix, g is the SM $\text{SU}(2)_L$ gauge coupling, M_W is the mass of the W boson given by the Higgs vacuum expectation value $gv/2$, and α, β, γ , and δ are flavor indices. The chiral couplings $g_{L,f}$ and $g_{R,f}$ are defined as

$$g_{L,f} = T_f^3 - Q_f \sin^2 \theta_W, \quad g_{R,f} = -Q_f \sin^2 \theta_W, \quad (2.3)$$

where Q_f is the fermion's charge in units of $+e$. We choose a flavor basis such that the down-type quark and the lepton Yukawa matrices are diagonal for convenience of calculation. The transformation for the up-type quarks between the flavor (primed) and mass basis (unprimed) reads

$$u'_{L,\alpha} = V_{\alpha\beta}^\dagger u_{L,\beta}. \quad (2.4)$$

In the spirit of EFT, the theory is valid only at low energies, $E \ll M_W$. Note that NSI [7–9] are of the same form as eqs. (2.1) and (2.2), but the scale and couplings are free parameters to reflect the unknown nature of new physics.

2.1 LEFT

Going beyond the SM, the full list of dimension-six four-fermion SMNEFT operators, which include the left-handed and right-handed neutrino states, are given in ref. [14]. Integrating out W^\pm, Z, h and t leads to the low-energy effective field theory (LEFT), respecting $\text{SU}(3)_C \times \text{U}(1)_Q$. GNI via neutral and charged currents, containing scalar, pseudoscalar, vector, axial vector and tensor terms at dimension-six level can be parameterized as

$$\mathcal{L}_{\text{LEFT}}^{\text{NC}} \supset -\frac{G_F}{\sqrt{2}}\sum_{j=1}^{10}(\tilde{\epsilon}_{j,f}^{(\sim)})^{\alpha\beta\gamma\delta}(\bar{\nu}_\alpha O_j \nu_\beta)(\bar{f}_\gamma O'_j f_\delta), \quad (2.5)$$

$$\mathcal{L}_{\text{LEFT}}^{\text{CC}} \supset -\frac{G_F V_{\delta\gamma}^*}{\sqrt{2}}\sum_{j=1}^{10}(\tilde{\epsilon}_{j,du}^{(\sim)})^{\alpha\beta\gamma\delta}(\bar{\nu}_\alpha O_j \ell_\beta)(\bar{d}_\gamma O'_j u_\delta) + \text{h.c.}, \quad (2.6)$$

where the operators O_j, O'_j and parameters $\tilde{\epsilon}^{(\sim)}$ are listed in table 1. The Dirac spinor $\nu_\alpha = (\nu_{L\alpha}, N_\alpha)^T$, and u and d indicate the mass eigenstates of up- and down-type quarks,

j	$(\tilde{\epsilon})_j$	O_j	O'_j
1	ϵ_L	$\gamma_\mu(\mathbb{1} - \gamma^5)$	$\gamma_\mu(\mathbb{1} - \gamma^5)$
2	$\tilde{\epsilon}_L$	$\gamma_\mu(\mathbb{1} + \gamma^5)$	$\gamma_\mu(\mathbb{1} - \gamma^5)$
3	ϵ_R	$\gamma_\mu(\mathbb{1} - \gamma^5)$	$\gamma_\mu(\mathbb{1} + \gamma^5)$
4	$\tilde{\epsilon}_R$	$\gamma_\mu(\mathbb{1} + \gamma^5)$	$\gamma_\mu(\mathbb{1} + \gamma^5)$
5	ϵ_S	$\mathbb{1} - \gamma^5$	$\mathbb{1}$
6	$\tilde{\epsilon}_S$	$\mathbb{1} + \gamma^5$	$\mathbb{1}$
7	$-\epsilon_P$	$\mathbb{1} - \gamma^5$	γ^5
8	$-\tilde{\epsilon}_P$	$\mathbb{1} + \gamma^5$	γ^5
9	ϵ_T	$\sigma_{\mu\nu}(\mathbb{1} - \gamma^5)$	$\sigma_{\mu\nu}(\mathbb{1} - \gamma^5)$
10	$\tilde{\epsilon}_T$	$\sigma_{\mu\nu}(\mathbb{1} + \gamma^5)$	$\sigma_{\mu\nu}(\mathbb{1} + \gamma^5)$

Table 1. Effective coupling constants and operators.

respectively.¹ Hermiticity of the Lagrangian requires the scalar and tensor effective couplings to satisfy

$$\begin{aligned}
 \epsilon_{S,f}^{\alpha\beta\gamma\delta} &= (\tilde{\epsilon}_{S,f}^{\beta\alpha\delta\gamma})^*, & \epsilon_{P,f}^{\alpha\beta\gamma\delta} &= -(\tilde{\epsilon}_{P,f}^{\beta\alpha\delta\gamma})^*, & \epsilon_{T,f}^{\alpha\beta\gamma\delta} &= (\tilde{\epsilon}_{T,f}^{\beta\alpha\delta\gamma})^*, \\
 \epsilon_{S,du}^{\alpha\beta\gamma\delta} &= (\tilde{\epsilon}_{S,du}^{\beta\alpha\delta\gamma})^*, & \epsilon_{P,du}^{\alpha\beta\gamma\delta} &= -(\tilde{\epsilon}_{P,du}^{\beta\alpha\delta\gamma})^*, & \epsilon_{T,du}^{\alpha\beta\gamma\delta} &= (\tilde{\epsilon}_{T,du}^{\beta\alpha\delta\gamma})^*.
 \end{aligned} \tag{2.7}$$

If the BSM new physics scale is Λ with a typical tree-level coupling κ , then parametrically $(\tilde{\epsilon}) \sim \kappa^2 v^2 / \Lambda^2$. Note that the operators with $j = 1, 3$ are the familiar NSI terms, and are a subset of SMEFT.

2.2 SMNEFT

Of the dim-6 4-fermion SMNEFT operators related to scalar and tensor GNI, the three chirality-flipping operators that couple to quarks are

1. $O_{NLQu}^{\alpha\beta\gamma\delta} = (\bar{N}_\alpha L_\beta^j)(\bar{Q}_\gamma^j u_\delta)$,
2. $O_{NLdQ}^{\alpha\beta\gamma\delta} = (\bar{N}_\alpha L_\beta^j)\epsilon_{jk}(\bar{d}_\gamma Q_\delta^k)$,
3. $O'_{NLdQ}^{\alpha\beta\gamma\delta} = (\bar{N}_\alpha \sigma_{\mu\nu} L_\beta^j)\epsilon_{jk}(\bar{d}_\gamma \sigma^{\mu\nu} Q_\delta^k)$,

where the fields are written in two-component spinors. L and Q are the left-handed lepton and quark doublet, respectively, and N is the right-handed neutrino state. Here, $\sigma^{\mu\nu} = \frac{i}{2}[\sigma^\mu \bar{\sigma}^\nu - \sigma^\nu \bar{\sigma}^\mu]$, with $\sigma^\mu = (\mathbb{1}, \vec{\sigma})$ and $\bar{\sigma}^\mu = (\mathbb{1}, -\vec{\sigma})$. We do not consider other dim-6 4-fermion SMNEFT operators since they lead to nonstandard charged lepton interactions and are therefore strongly constrained [14]. We can write the effective Lagrangian as

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + 2\sqrt{2}G_F[C_{NLdQ}O_{NLdQ} + C_{NLQu}O_{NLQu} + C'_{NLdQ}O'_{NLdQ}], \tag{2.8}$$

¹Our analysis can also be applied to Majorana neutrinos with the neutrino bilinears in eqs. (2.1) and (2.5) replaced by $\bar{N}_{M\alpha}O\nu_{M\beta}$, where the Majorana spinors are $\nu_M = (\nu_L, \nu_L^c)^T$, $N_M = (N^c, N)^T$.

where the flavor indices are omitted for simplicity. In the same spirit of power counting as in the last subsection, the Wilson coefficients (WCs) have the general dependence $C \sim \kappa^2 v^2 / \Lambda^2$. For instance, $C \sim \mathcal{O}(10^{-4})$ if $\Lambda \sim 10 \text{ TeV}$ and $\kappa \sim 1$.

To jointly interpret the results of experiments at very different energy scales, a consistent theoretical framework is needed. LEFT and SMNEFT are the language we use to describe the physics below and above the electroweak scale v , respectively. The renormalization group (RG) running below and above the electroweak scale makes it possible to directly compare low-energy and high-energy probes. Leading-order (LO) matching between these two EFTs is performed at the electroweak scale.

2.3 RG running and matching

Since we will use both low-energy neutrino scattering experiments and high-energy colliders to constrain these Wilson coefficients (WCs), renormalization group (RG) running and matching have to be implemented. We perform leading-order (LO) matching of these two EFTs at the electroweak scale:

$$\begin{aligned}
 \epsilon_{S,d}^{\alpha\beta\gamma\delta} &= -C_{NLdQ}^{\alpha\beta\gamma\delta}, & \epsilon_{S,u}^{\alpha\beta\gamma\delta} &= -C_{NLQu}^{\alpha\beta\rho\delta} V_{\rho\gamma}, \\
 \epsilon_{P,d}^{\alpha\beta\gamma\delta} &= -C_{NLdQ}^{\alpha\beta\gamma\delta}, & \epsilon_{P,u}^{\alpha\beta\gamma\delta} &= C_{NLQu}^{\alpha\beta\rho\delta} V_{\rho\gamma}, \\
 \epsilon_{T,d}^{\alpha\beta\gamma\delta} &= -C_{NLdQ}^{\prime\alpha\beta\gamma\delta}, \\
 \epsilon_{S,du}^{\alpha\beta\gamma\delta} &= \frac{C_{NLdQ}^{\alpha\beta\gamma\rho} V_{\rho\delta}^\dagger - C_{NLQu}^{\alpha\beta\gamma\delta}}{V_{\delta\gamma}^*}, & \epsilon_{P,du}^{\alpha\beta\gamma\delta} &= \frac{C_{NLdQ}^{\alpha\beta\gamma\rho} V_{\rho\delta}^\dagger + C_{NLQu}^{\alpha\beta\gamma\delta}}{V_{\delta\gamma}^*}, \\
 \epsilon_{T,du}^{\alpha\beta\gamma\delta} &= C_{NLdQ}^{\prime\alpha\beta\gamma\rho} \frac{V_{\rho\delta}^\dagger}{V_{\delta\gamma}^*}.
 \end{aligned} \tag{2.9}$$

As we run down, both neutral and charged current WCs are induced by each of the three SMNEFT operators. Therefore they are not independent of each other. Their relations at the electroweak scale are

$$\begin{aligned}
 \epsilon_{S,d}^{\alpha\beta\gamma\delta} &= -\frac{V_{\gamma\rho}^\dagger}{2V_{\delta\rho}} (\epsilon_{S,du}^{\alpha\beta\gamma\rho} + \epsilon_{P,du}^{\alpha\beta\gamma\rho}), & \epsilon_{P,d}^{\alpha\beta\gamma\delta} &= -\epsilon_{S,d}^{\alpha\beta\gamma\delta}, \\
 \epsilon_{S,u}^{\alpha\beta\gamma\delta} &= \frac{1}{2} V_{\gamma\rho} V_{\delta\rho}^* (\epsilon_{S,du}^{\alpha\beta\gamma\rho} - \epsilon_{P,du}^{\alpha\beta\gamma\rho}), & \epsilon_{P,u}^{\alpha\beta\gamma\delta} &= -\epsilon_{S,u}^{\alpha\beta\gamma\delta}, \\
 \epsilon_{T,d}^{\alpha\beta\gamma\delta} &= -\epsilon_{T,du}^{\alpha\beta\gamma\rho} \frac{V_{\rho\gamma}}{V_{\rho\delta}^*}.
 \end{aligned} \tag{2.10}$$

We have performed the RG running above and below the weak scale, the details of which are described in appendix A. The RG equations are run from 2 GeV to 1 TeV, which is the typical LHC scale. Eventually we place bounds on the SMNEFT WCs at 1 TeV. The

anomalous dimension matrix we calculated at the one-loop level is

$$\mu \frac{d}{d\mu} \begin{pmatrix} C_{NLQu} \\ C_{NLdQ} \\ C'_{NLdQ} \end{pmatrix}_{(\mu)} = \left[\frac{\alpha_1(\mu)}{2\pi} \begin{pmatrix} -1/3 & 0 & 0 \\ 0 & 1/6 & -1 \\ 0 & -1/48 & -5/9 \end{pmatrix} + \frac{\alpha_2(\mu)}{2\pi} \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 9 \\ 0 & 3/16 & -3/2 \end{pmatrix} \right. \\ \left. + \frac{\alpha_3(\mu)}{2\pi} \begin{pmatrix} -4 & 0 & 0 \\ 0 & -4 & 0 \\ 0 & 0 & 4/3 \end{pmatrix} \right] \begin{pmatrix} C_{NLQu} \\ C_{NLdQ} \\ C'_{NLdQ} \end{pmatrix}_{(\mu)}, \quad (2.11)$$

$$\mu \frac{d}{d\mu} \begin{pmatrix} \epsilon_{S,du} \\ \epsilon_{P,du} \\ \epsilon_{T,du} \end{pmatrix}_{(\mu)} = \left[\frac{\alpha_e(\mu)}{2\pi} \begin{pmatrix} 2/3 & 0 & 4 \\ 0 & 2/3 & 4 \\ 1/24 & 1/24 & -20/9 \end{pmatrix} \right. \\ \left. + \frac{\alpha_3(\mu)}{2\pi} \begin{pmatrix} -4 & 0 & 0 \\ 0 & -4 & 0 \\ 0 & 0 & 4/3 \end{pmatrix} \right] \begin{pmatrix} \epsilon_{S,du} \\ \epsilon_{P,du} \\ \epsilon_{T,du} \end{pmatrix}_{(\mu)}, \quad (2.12)$$

$$\mu \frac{d}{d\mu} \begin{pmatrix} \epsilon_{S,d} \\ \epsilon_{P,d} \\ \epsilon_{T,d} \end{pmatrix}_{(\mu)} = \left[\frac{\alpha_e(\mu)}{2\pi} \begin{pmatrix} -1/9 & 0 & 0 \\ 0 & -1/9 & 0 \\ 0 & 0 & 5/36 \end{pmatrix} \right. \\ \left. + \frac{\alpha_3(\mu)}{2\pi} \begin{pmatrix} -4 & 0 & 0 \\ 0 & -4 & 0 \\ 0 & 0 & 4/3 \end{pmatrix} \right] \begin{pmatrix} \epsilon_{S,d} \\ \epsilon_{P,d} \\ \epsilon_{T,d} \end{pmatrix}_{(\mu)}, \quad (2.13)$$

$$\mu \frac{d}{d\mu} \begin{pmatrix} \epsilon_{S,u} \\ \epsilon_{P,u} \\ \epsilon_{T,u} \end{pmatrix}_{(\mu)} = \left[\frac{\alpha_e(\mu)}{2\pi} \begin{pmatrix} -4/9 & 0 & 0 \\ 0 & -4/9 & 0 \\ 0 & 0 & 5/9 \end{pmatrix} \right. \\ \left. + \frac{\alpha_3(\mu)}{2\pi} \begin{pmatrix} -4 & 0 & 0 \\ 0 & -4 & 0 \\ 0 & 0 & 4/3 \end{pmatrix} \right] \begin{pmatrix} \epsilon_{S,u} \\ \epsilon_{P,u} \\ \epsilon_{T,u} \end{pmatrix}_{(\mu)}, \quad (2.14)$$

where the flavor indices are implicit. The QED and weak couplings are important as they introduce mixing between different operators. Solving the differential equations with the three-loop β -functions and taking into account the top and bottom quark mass thresholds, we obtain the numerical relations between effective couplings at different energy scales:

$$\begin{pmatrix} C_{NLQu} \\ C_{NLdQ} \\ C'_{NLdQ} \end{pmatrix}_{(\mu=M_Z)} = \begin{pmatrix} 1.18 & 0 & 0 \\ 0 & 1.18 & -0.117 \\ 0 & -2.44 \times 10^{-3} & 0.966 \end{pmatrix} \begin{pmatrix} C_{NLQu} \\ C_{NLdQ} \\ C'_{NLdQ} \end{pmatrix}_{(\mu=1 \text{ TeV})}, \quad (2.15)$$

$$\begin{pmatrix} \epsilon_{S,du} \\ \epsilon_{P,du} \\ \epsilon_{T,du} \end{pmatrix}_{(\mu=2\text{ GeV})} = \begin{pmatrix} 1.52 & 2.34 \times 10^{-6} & -0.0218 \\ 2.34 \times 10^{-6} & 1.52 & -0.0218 \\ -2.26 \times 10^{-4} & -2.26 \times 10^{-4} & 0.878 \end{pmatrix} \begin{pmatrix} \epsilon_{S,du} \\ \epsilon_{P,du} \\ \epsilon_{T,du} \end{pmatrix}_{(\mu=M_Z)}, \quad (2.16)$$

$$\begin{pmatrix} \epsilon_{S,d} \\ \epsilon_{P,d} \\ \epsilon_{T,d} \end{pmatrix}_{(\mu=2\text{ GeV})} = \begin{pmatrix} 1.52 & 0 & 0 \\ 0 & 1.52 & 0 \\ 0 & 0 & 0.869 \end{pmatrix} \begin{pmatrix} \epsilon_{S,d} \\ \epsilon_{P,d} \\ \epsilon_{T,d} \end{pmatrix}_{(\mu=M_Z)}, \quad (2.17)$$

$$\begin{pmatrix} \epsilon_{S,u} \\ \epsilon_{P,u} \\ \epsilon_{T,u} \end{pmatrix}_{(\mu=2\text{ GeV})} = \begin{pmatrix} 1.53 & 0 & 0 \\ 0 & 1.53 & 0 \\ 0 & 0 & 0.867 \end{pmatrix} \begin{pmatrix} \epsilon_{S,u} \\ \epsilon_{P,u} \\ \epsilon_{T,u} \end{pmatrix}_{(\mu=M_Z)}. \quad (2.18)$$

The numerical relations between LEFT WCs at 2 GeV and SMNEFT WCs at 1 TeV, with $V_{ud} = 0.97420$ [21], are

$$\begin{aligned} \epsilon_{S,du} &= -1.84C_{NLQu} + 1.79C_{NLdQ} - 0.199C'_{NLdQ}, \\ \epsilon_{P,du} &= 1.84C_{NLQu} + 1.79C_{NLdQ} - 0.157C'_{NLdQ}, \\ \epsilon_{T,du} &= 5.49 \times 10^{-4}C_{NLQu} - 2.14 \times 10^{-3}C_{NLdQ} + 0.849C'_{NLdQ}, \\ \epsilon_{S,u} &= -1.76C_{NLQu}, \\ \epsilon_{P,u} &= 1.76C_{NLQu}, \\ \epsilon_{T,u} &= 0, \\ \epsilon_{S,d} &= -1.80C_{NLdQ} + 0.179C'_{NLdQ}, \\ \epsilon_{P,d} &= -1.80C_{NLdQ} + 0.179C'_{NLdQ}, \\ \epsilon_{T,d} &= 2.12 \times 10^{-3}C_{NLdQ} - 0.839C'_{NLdQ}. \end{aligned} \quad (2.19)$$

Low energy constraints on the SMNEFT WCs from nuclear beta decay, pseudoscalar meson decay, and coherent scattering have been discussed in ref. [14] without accounting for the effects of RG running. The RG running is crucial, as it introduces operator mixing which produces degeneracies in the WCs. Here we first calculate the LEFT and SMNEFT WCs below and above the electroweak scale, respectively. After the RG running, we convert the low energy constraints on the LEFT WCs to the high energy constraints on the SMNEFT WCs, and compare them with those from high energy collider experiments at the same energy scale.

3 Low-energy constraints

3.1 Neutrino mass bounds

Scalar and tensor interactions that flip the neutrino chirality contribute to the neutrino mass radiatively. Both one- and two-loop corrections to the neutrino mass can be generated by chirality-changing operators. Here we ignore the one-loop corrections since, except for the top quark, they are (counterintuitively) suppressed by a factor of $(m_q/M_Z)^2$ as

compared to the two-loop corrections [22, 23]. The two-loop contribution is estimated as

$$\Delta m_\nu \simeq 3g^2 G_F \epsilon \frac{m_q M_W^2}{(4\pi)^4} \left(\ln \frac{\mu^2}{M_W^2} \right)^2, \quad (3.1)$$

where m_q is a quark mass, μ is the renormalization scale, and ϵ can be either a NC or CC GNI parameter. We conservatively take μ to not be too far above the electroweak scale so that the top quark loop correction is suppressed.

Bounds from neutrino masses and oscillations are very model specific because of the importance of the properties of the particles in the loops and the possibility of cancellations between loop and other contributions. However, barring fine-tuned cancellations, they provide an order of magnitude estimate of how much the new interactions may contribute to neutrino masses. For our estimates, we assume neutrinos acquire mass only from loop effects due to the new interactions, i.e., neutrino masses vanish as $\epsilon \rightarrow 0$. Then, constraints on the contact interactions can be obtained by requiring $\Delta m_\nu < \sum m_\nu$. A recent upper bound on the sum of neutrino mass from cosmological observations and particle physics experiments is $\sum m_\nu \lesssim 0.26$ eV [24], which is model dependent. The most recent model-independent bound is that obtained by the KATRIN Collaboration [25]. They reported a 1.1 eV upper bound on the effective neutrino mass based on the β -decay electron spectrum. The bounds on the scalar and tensor contact interactions from neutrino masses without (with) cosmological inputs are

$$|\epsilon_{S,P,T}^{\alpha\beta 11}| \lesssim 10^{-3} (10^{-4}), \quad |\epsilon_{S,P,T}^{\alpha\beta 22}| \lesssim 10^{-5} (10^{-6}), \quad |\epsilon_{S,P,T}^{\alpha\beta 33}| \lesssim 10^{-6} (10^{-7}). \quad (3.2)$$

The bounds using cosmological data are only suggestive because we have not evaluated how the relic neutrino abundance is affected by the new interactions. From eq. (3.2), we see that if GNI are also coupled to heavy quark flavors, the bounds on the SMNEFT WCs C_{NLQu} , C_{NLdQ} , and C'_{NLdQ} are too strong to be probed by other experiments, current or future. Despite the highly model-dependent nature of this conclusion, we focus on couplings to first generation quarks in the rest of the paper.

Related bounds arise from neutrino magnetic moments via an external photon attached to the fermion loop responsible for neutrino mass generation. The magnetic moment induced by scalar and tensor GNI is bounded by [26]

$$\mu_\nu \approx \frac{e G_F m_d}{8\pi^2} \epsilon \lesssim 3 \times 10^{-11} \mu_B, \quad (3.3)$$

where the Bohr magneton $\mu_B = \frac{e\hbar}{2m_e c} \simeq 2.9 \times 10^{-7}$ eV⁻¹. This yields

$$|\epsilon_{S,P,T}^{\alpha\beta 11}| \lesssim 30, \quad (3.4)$$

which are much weaker than the bounds above.

3.2 Pseudoscalar meson decay

The pseudoscalar quark bilinear can contribute to the leptonic decay of a pseudoscalar meson (P). In the SM, the decay is helicity suppressed so that the width $\Gamma_{\text{SM}}(P \rightarrow \ell\nu) \propto$

m_ℓ^2 . The suppression is lifted by pseudoscalar GNI

$$\Gamma_{\text{GNI,p}}(P \rightarrow \ell \beta \nu_\alpha) \propto (\epsilon_{P,du}^{\alpha\beta 11})^2 \frac{m_\pi^4}{(m_u + m_d)^2}. \quad (3.5)$$

The branching ratio

$$R_\pi \equiv \frac{\Gamma(\pi \rightarrow e \nu[\gamma])}{\Gamma(\pi \rightarrow \mu \nu[\gamma])} = R_\pi^{(0)}[1 + \Delta_\pi], \quad \text{with} \quad R_\pi^{(0)} = \frac{m_e^2}{m_\mu^2} \left(\frac{m_\pi^2 - m_e^2}{m_\pi^2 - m_\mu^2} \right)^2, \quad (3.6)$$

serves as a good observable, as the experiment systematic uncertainties shared by the two processes cancel in the ratio. Δ_π contains higher order corrections [27]. $\Gamma_{(\pi \rightarrow \ell \bar{\nu}[\gamma])}$ contains physical and virtual photons (radiative corrections). Including pseudoscalar GNI interactions [28],

$$\frac{R_\pi}{R_\pi^{\text{SM}}} = \frac{1 + |\frac{B_0}{m_e} \epsilon_{P,du}^{\alpha e 11}|^2}{1 + |\frac{B_0}{m_\mu} \epsilon_{P,du}^{\alpha \mu 11}|^2}, \quad (3.7)$$

where $B_0(\mu) = m_\pi^2/(m_u(\mu) + m_d(\mu))$. Taking $m_\pi = 139.57 \text{ MeV}$, $m_u^{\overline{\text{MS}}}(\mu = 2 \text{ GeV}) = 2.16 \text{ MeV}$ and $m_d^{\overline{\text{MS}}}(\mu = 2 \text{ GeV}) = 4.67 \text{ MeV}$ [21], gives $B_0^{\overline{\text{MS}}}(\mu = 2 \text{ GeV}) = 2.8 \times 10^3 \text{ MeV}$. The current combined uncertainty in R_π^{exp} [21, 29–31] and R_π^{SM} [27, 32] are

$$R_\pi = 1.2327(23) \times 10^{-4}, \quad R_\pi^{\text{SM}} = 1.2352(1) \times 10^{-4}. \quad (3.8)$$

If both $\epsilon_{P,du}^{\alpha e 11}$ and $\epsilon_{P,du}^{\alpha \mu 11}$ are allowed to vary simultaneously, no bound on either parameter is obtained because they are degenerate, as is evident from eq. (3.7). With the assumption that only one of $\epsilon_{P,du}$ is nonzero, the 90% C.L. bounds are

$$|\epsilon_{P,du}^{\alpha e 11}| < 6.2 \times 10^{-6}, \quad \text{and} \quad |\epsilon_{P,du}^{\alpha \mu 11}| < 2.7 \times 10^{-3}. \quad (3.9)$$

Because the measured branching to the electron channel is tiny, $\epsilon_{P,du}^{\alpha e 11}$ is highly constrained. These bounds are much stronger than the ones obtained in ref. [28], which assumed that both $\epsilon_{P,du}$ and $\tilde{\epsilon}_{P,du}$ are simultaneously nonzero, which however, cannot be realized with the three SMNEFT operators considered here. The bounds on the coefficients of the low-energy effective Lagrangian can be translated to bounds on the three SMNEFT WCs by adopting the relations in eq. (2.19), which display degeneracies between the SMNEFT WCs. We therefore bound the individual WCs by setting the other two to zero. The 90% C.L. bounds on the SMNEFT WCs are

$$|C_{NLQu}^{\alpha e 11}| < 3.3 \times 10^{-6}, \quad |C_{NLdQ}^{\alpha e 11}| < 3.4 \times 10^{-6}, \quad |C_{NLdQ}'^{\alpha e 11}| < 3.9 \times 10^{-5}, \quad (3.10)$$

$$|C_{NLQu}^{\alpha \mu 11}| < 1.5 \times 10^{-3}, \quad |C_{NLdQ}^{\alpha \mu 11}| < 1.5 \times 10^{-3}, \quad |C_{NLQu}'^{\alpha \mu 11}| < 1.7 \times 10^{-2}. \quad (3.11)$$

The correlations between the C_{NLdQ} and C_{NLQu} (C_{NLdQ}' and C_{NLQu}'), with C_{NLdQ}' (C_{NLQu}') set to zero, are shown by the green lines in the upper (lower) panel of figure 1.

To circumvent the degeneracy in $\epsilon_{P,du}^{\alpha e 11}$ and $\epsilon_{P,du}^{\alpha \mu 11}$ in eq. (3.7), we now apply the individual decay width measurements of $\pi \rightarrow \ell \nu[\gamma]$ to set the bounds. In the SM, the decay width at tree level is

$$\Gamma_{(\pi \rightarrow \ell \bar{\nu}[\gamma])} = \frac{G_F^2}{8\pi} f_\pi^2 V_{ud}^2 m_\ell^2 m_\pi \left(1 - \frac{m_\ell^2}{m_\pi^2} \right)^2 (1 + \Delta_\pi). \quad (3.12)$$

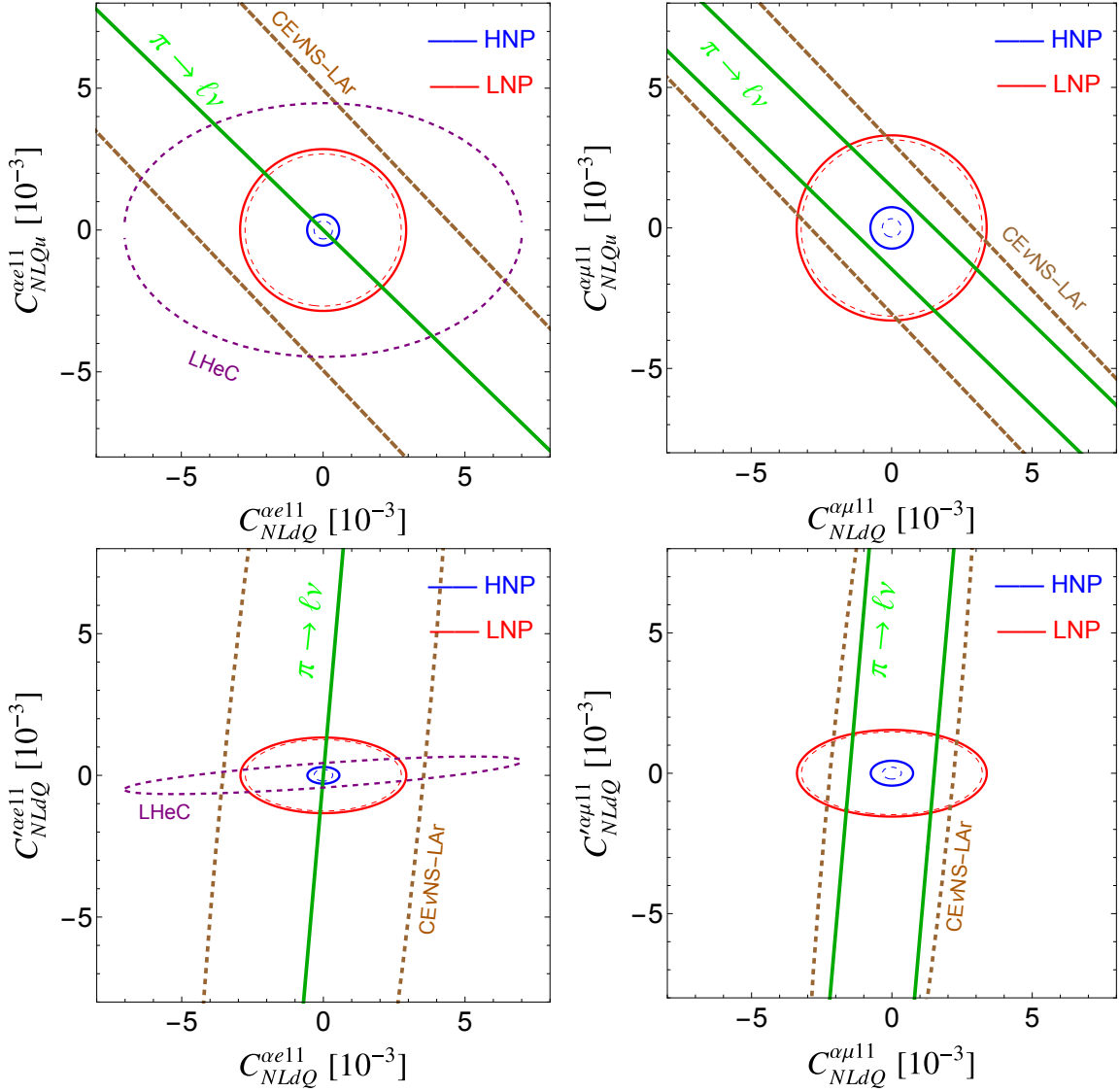


Figure 1. The 90% C.L. allowed regions in the C_{NLdQ} - C_{NLQu} planes (upper panels) and C_{NLdQ} - C'_{NLdQ} planes (lower panels) at 1 TeV with electron flavor (left panels) and muon flavor (right panels). The green lines (overlapping in the left panels) are the bounds from pion decay with the third parameter is set to zero to break the degeneracy. The red (blue) solid contours correspond to current LHC searches with $L = 139 \text{ fb}^{-1}$ for the low-scale new physics LNP (high-scale new physics HNP) case. The brown dashed lines show the projected bounds from the future LAr COHERENT experiment, with C_{NLQu} is set to zero in the lower panels to obtain meaningful bounds. The red and blue dashed contours are the projected bounds from HL-LHC with 3 ab^{-1} of data. The dashed purple contours in the left panels are the projected bounds from LHeC with 3 ab^{-1} .

The theoretical uncertainties are mainly from calculations of the decay constant and radiative corrections. According to refs. [21, 32–35],

$$\Delta_\pi = 0.0176 \pm 0.0021, \quad f_\pi = 130.2 \pm 1.2 \text{ MeV}. \quad (3.13)$$

The universal theoretical uncertainties yield

$$\frac{\delta\Gamma(\pi^+ \rightarrow e^+ \nu_e [\gamma])}{\Gamma(\pi^+ \rightarrow e^+ \nu_e [\gamma])} = \frac{\delta\Gamma(\pi^+ \rightarrow \mu^+ \nu_\mu [\gamma])}{\Gamma(\pi^+ \rightarrow \mu^+ \nu_\mu [\gamma])} = 1.9 \times 10^{-2}. \quad (3.14)$$

Measurements give $BR(\pi^+ \rightarrow e^+ \nu_e [\gamma]) = (1.230 \pm 0.004) \times 10^{-4}\%$, $BR(\pi^+ \rightarrow \mu^+ \nu_\mu [\gamma]) = (99.98770 \pm 0.00004)\%$, and $\tau_{\pi^\pm} = 26.033(5) \text{ ns}$ [21]. The experimental uncertainties in the electron (muon) channel is 3.3×10^{-3} (1.9×10^{-4}) and can be neglected. Assuming that the new physics contributions do not exceed the theoretical uncertainties, the bounds on $\epsilon_{P,du}^{\alpha e 11}$ and $\epsilon_{P,du}^{\alpha \mu 11}$ are given by

$$\left| \frac{B_0}{m_e} \epsilon_{P,du}^{\alpha e 11} \right|^2 < \frac{\delta\Gamma(\pi^+ \rightarrow e^+ \nu_e [\gamma])}{\Gamma(\pi^+ \rightarrow e^+ \nu_e [\gamma])}, \quad \left| \frac{B_0}{m_\mu} \epsilon_{P,du}^{\alpha \mu 11} \right|^2 < \frac{\delta\Gamma(\pi^+ \rightarrow \mu^+ \nu_\mu [\gamma])}{\Gamma(\pi^+ \rightarrow \mu^+ \nu_\mu [\gamma])}, \quad (3.15)$$

which yield

$$|\epsilon_{P,du}^{\alpha e 11}| < 3.4 \times 10^{-5}, \quad |\epsilon_{P,du}^{\alpha \mu 11}| < 6.6 \times 10^{-3}, \quad (3.16)$$

at the 90% C.L. By allowing only one WC to be nonzero at a time, the 90% C.L. bounds on the SMNEFT WCs derived from individual decay channels are

$$|C_{NLQu}^{\alpha e 11}| < 1.9 \times 10^{-5}, \quad |C_{NLdQ}^{\alpha e 11}| < 1.9 \times 10^{-5}, \quad |C'_{NLdQ}{}^{\alpha e 11}| < 2.2 \times 10^{-4}, \quad (3.17)$$

$$|C_{NLQu}^{\alpha \mu 11}| < 3.6 \times 10^{-3}, \quad |C_{NLdQ}^{\alpha \mu 11}| < 3.7 \times 10^{-3}, \quad |C'_{NLQu}{}^{\alpha \mu 11}| < 4.2 \times 10^{-2}. \quad (3.18)$$

3.3 Nuclear beta decay

Nuclear β -decay is another low-energy probe that is sensitive to the new CC GNI interactions. The nucleon-level effective Lagrangian contributing to neutron beta decay, $n \rightarrow p + e^- + \bar{\nu}_e$, is using eq. (2.7),

$$\begin{aligned} \mathcal{L}_N = & -\frac{G_F}{\sqrt{2}} V_{ud} [\bar{p} \gamma^\mu (g_V - g_A \gamma^5) n \cdot \bar{e} \gamma_\mu (1 - \gamma^5) \nu_e + g_S (\epsilon_{S,du}^{\alpha e 11})^* \bar{p} n \cdot \bar{e} (1 + \gamma^5) \nu_e \\ & - g_P (\epsilon_{P,du}^{\alpha e 11})^* \bar{p} \gamma^5 n \cdot \bar{e} (1 + \gamma^5) \nu_e + 2g_T (\epsilon_{T,du}^{\alpha e 11})^* \bar{p} \sigma^{\mu\nu} n \cdot \bar{e} \sigma_{\mu\nu} (1 + \gamma^5) \nu_e] + \text{h.c.}, \end{aligned} \quad (3.19)$$

where $g_{V(A)}$ is the (axial-)vector charge and $g_{S,P,T}$ are the nonstandard charges. Neglecting nucleon recoil and the pseudoscalar contribution in the $q^2 \rightarrow 0$ limit, the neutron β decay width is

$$\Gamma = \frac{G_F^2 V_{ud}^2}{2\pi^3} [g_V^2 (3\lambda^2 + 1) + g_S^2 |\epsilon_{S,du}^{\alpha e 11}|^2 + 48g_T^2 |\epsilon_{T,du}^{\alpha e 11}|^2] I, \quad (3.20)$$

where $\lambda \equiv g_V/g_A$ and

$$I = \int p_e E_e (M_n - M_p - E_e)^2 dE_e \approx 0.06 \text{ MeV}^5. \quad (3.21)$$

The decay width can also be written in terms of the NC effective couplings by using the relations in eq. (2.10):

$$\Gamma = \frac{G_F^2 V_{ud}^2}{2\pi^3} [g_V^2 (3\lambda^2 + 1) + g_S^2 \left(\frac{\epsilon_{S,u}^{\alpha e 11}}{V_{ud}^2} - \epsilon_{S,d}^{\alpha e 11} \right)^2 + 48g_T^2 (\epsilon_{T,d}^{\alpha e 11})^2] I. \quad (3.22)$$

From ref. [36], the 90% C.L. bounds, based on the differential observables from polarized nuclear beta decay, are

$$|\epsilon_{S,du}^{\alpha e 11}| < 0.063, \quad |\epsilon_{T,du}^{\alpha e 11}| < 0.024. \quad (3.23)$$

Bounds on the NC parameters can be computed by using the relations in eq. (2.10) with $\epsilon_{P,du}^{\alpha e 11}$ taken to be 0:

$$|\epsilon_{S,d}^{\alpha e 11}|, |\epsilon_{P,d}^{\alpha e 11}| < 0.063, \quad |\epsilon_{S,u}^{\alpha e 11}|, |\epsilon_{P,u}^{\alpha e 11}| < 0.060, \quad |\epsilon_{T,d}^{\alpha e 11}| < 0.024. \quad (3.24)$$

Degeneracies do not permit simultaneous bounds on all the SMNEFT WCs. With the assumption that only one of them is nonzero, the 90% C.L. bounds are

$$|C_{NLQu}^{\alpha e 11}| < 3.4 \times 10^{-2}, \quad |C_{NLdQ}^{\alpha e 11}| < 3.5 \times 10^{-2}, \quad |C_{NLdQ}'^{\alpha e 11}| < 2.8 \times 10^{-2}. \quad (3.25)$$

These constraints are much weaker than the ones from charged pion decay.

3.4 Neutrino deep inelastic scattering

Neutrino deep inelastic scattering on nucleons can be modified by scalar, pseudoscalar, and tensor GNI. Please note that the charged current cannot be affected by the three GNI considered in this paper, as the right-handed neutrino is absent in the neutrino beams. The total charged current and neutral current neutrino-nucleon scattering cross sections in the SM are

$$\sigma_{\nu N,SM}^{CC} = \frac{2G_F^2}{\pi} E_\nu M_N \left[\langle x d_N + x \frac{1}{3} \bar{u}_N \rangle \right], \quad (3.26)$$

$$\sigma_{\bar{\nu} N,SM}^{CC} = \frac{2G_F^2}{\pi} E_\nu M_N \left[\langle x \frac{1}{3} u_N + x \bar{d}_N \rangle \right], \quad (3.27)$$

$$\begin{aligned} \sigma_{\nu N,SM}^{NC} = \frac{2G_F^2}{\pi} E_\nu M_N & \left[\left(g_{L,u}^2 + \frac{1}{3} g_{R,u}^2 \right) \langle x u_N \rangle + \left(g_{L,d}^2 + \frac{1}{3} g_{R,d}^2 \right) \langle x d_N \rangle \right. \\ & \left. + \left(g_{R,u}^2 + \frac{1}{3} g_{L,u}^2 \right) \langle x \bar{u}_N \rangle + \left(g_{R,d}^2 + \frac{1}{3} g_{L,d}^2 \right) \langle x \bar{d}_N \rangle \right], \end{aligned} \quad (3.28)$$

$$\begin{aligned} \sigma_{\bar{\nu} N,SM}^{NC} = \frac{2G_F^2}{\pi} E_\nu M_N & \left[\left(g_{R,u}^2 + \frac{1}{3} g_{L,u}^2 \right) \langle x u_N \rangle + \left(g_{R,d}^2 + \frac{1}{3} g_{L,d}^2 \right) \langle x d_N \rangle \right. \\ & \left. + \left(g_{L,u}^2 + \frac{1}{3} g_{R,u}^2 \right) \langle x \bar{u}_N \rangle + \left(g_{L,d}^2 + \frac{1}{3} g_{R,d}^2 \right) \langle x \bar{d}_N \rangle \right], \end{aligned} \quad (3.29)$$

where we have neglected contributions from heavy quarks, and

$$\langle x q_N \rangle \equiv \int_0^1 x q_N(x) dx, \quad \langle x \bar{q}_N \rangle \equiv \int_0^1 x \bar{q}_N(x) dx, \quad (3.30)$$

determine the fraction of nucleon momentum carried by quarks and anti-quarks. $g_{L,f}$ and $g_{R,f}$ are the SM effective couplings given in eq. (2.3). We take [37]

$$g_{L,u} = 0.3457, \quad g_{R,u} = -0.1553, \quad g_{L,d} = -0.4288, \quad g_{R,d} = 0.0777, \quad (3.31)$$

which include the one-loop and leading two-loop corrections. The neutral current is modified by scalar, pseudoscalar, and tensor GNI:

$$\sigma_{\nu N, S(P)}^{NC} = \sigma_{\bar{\nu} N, S(P)}^{NC} = \frac{G_F^2}{12\pi} E_\nu M_N [\epsilon_{S(P),u}^2 \langle x(u_N + \bar{u}_N) \rangle + \epsilon_{S(P),d}^2 \langle x(d_N + \bar{d}_N) \rangle], \quad (3.32)$$

$$\sigma_{\nu N, T}^{NC} = \sigma_{\bar{\nu} N, T}^{NC} = \frac{56G_F^2}{3\pi} E_\nu M_N [\epsilon_{T,u}^2 \langle x(u_N + \bar{u}_N) \rangle + \epsilon_{T,d}^2 \langle x(d_N + \bar{d}_N) \rangle], \quad (3.33)$$

where the flavor indices are suppressed for simplicity. In the following analysis, we assume the target is isoscalar and composed of free nucleons, so that we may use the proton PDF. Under these assumptions, the nuclear PDFs become

$$\langle x d_N \rangle = \langle x u_N \rangle = \frac{N}{2} \langle x(u_p + d_p) \rangle, \quad \langle x \bar{d}_N \rangle = \langle x \bar{u}_N \rangle = \frac{N}{2} \langle x(\bar{d}_p + \bar{u}_p) \rangle. \quad (3.34)$$

3.4.1 CHARM: $\nu_e q \rightarrow \nu q$

The CHARM collaboration measured the ratio of total cross sections for semileptonic ν_e and $\bar{\nu}_e$ scattering to be [38]

$$R^e \equiv \frac{\sigma(\nu_e N \rightarrow \nu X) + \sigma(\bar{\nu}_e N \rightarrow \nu X)}{\sigma(\nu_e N \rightarrow e^- X) + \sigma(\bar{\nu}_e N \rightarrow e^+ X)} = 0.406 \pm 0.140. \quad (3.35)$$

The SM prediction from eqs. (3.26) to (3.29) is

$$R^e = g_L^2 + g_R^2 = 0.3335, \quad (3.36)$$

where

$$g_L^2 = g_{L,u}^2 + g_{L,d}^2, \quad g_R^2 = g_{R,u}^2 + g_{R,d}^2. \quad (3.37)$$

Including the new GNI contributions from eqs. (3.32) and (3.33), R^e becomes

$$R^e = g_L^2 + g_R^2 + \frac{1}{12} \sum_{q=u,d} ((\epsilon_{s,q}^{\alpha e 11})^2 + (\epsilon_{p,q}^{\alpha e 11})^2 + 224(\epsilon_{T,q}^{\alpha e 11})^2). \quad (3.38)$$

The 90% C.L. bounds on the LEFT parameters are

$$|\epsilon_{S,q}^{\alpha e 11}|, |\epsilon_{P,q}^{\alpha e 11}| < 1.9, \quad |\epsilon_{T,q}^{\alpha e 11}| < 0.13. \quad (3.39)$$

With only a single constraint on R^e , the degeneracy between the three SMNEFT WCs remains unbroken. The bounds on the SMNEFT WCs, with the assumption that only one of the WCs is nonzero at a time, are

$$|C_{NLQ_u}^{\alpha e 11}| < 0.77, \quad |C_{NLdQ}^{\alpha e 11}| < 0.75, \quad |C_{NLdQ}^{\prime \alpha e 11}| < 0.15, \quad (3.40)$$

which are much weaker than the bounds from charged pion decay and nuclear beta decay.

3.4.2 NuTeV: $\nu_\mu q \rightarrow \nu q$

The NuTeV collaboration has measured the ratios of neutral current to charged current neutrino-nucleon cross sections [39]:

$$R^\nu \equiv \frac{\sigma(\nu_\mu N \rightarrow \nu X)}{\sigma(\nu N \rightarrow \mu^- X)} = 0.3916 \pm 0.0013, \quad R^{\bar{\nu}} \equiv \frac{\sigma(\bar{\nu}_\mu N \rightarrow \bar{\nu} X)}{\sigma(\bar{\nu} N \rightarrow \mu^+ X)} = 0.4050 \pm 0.0027. \quad (3.41)$$

In the SM, the cross section ratios on an isoscalar target composed of free nucleons are

$$R_{SM}^\nu = \frac{(g_L^2 + \frac{1}{3}g_R^2)f_q + (g_R^2 + \frac{1}{3}g_L^2)f_{\bar{q}}}{f_q + \frac{1}{3}f_{\bar{q}}}, \quad R_{SM}^{\bar{\nu}} = \frac{(g_R^2 + \frac{1}{3}g_L^2)f_q + (g_L^2 + \frac{1}{3}g_R^2)f_{\bar{q}}}{\frac{1}{3}f_q + f_{\bar{q}}}, \quad (3.42)$$

where f_q and $f_{\bar{q}}$ determine the fraction of proton momentum carried by the first generation of quarks and anti-quarks:

$$f_q = \langle xu + xd \rangle = 0.42, \quad f_{\bar{q}} = \langle x\bar{u} + x\bar{d} \rangle = 0.068. \quad (3.43)$$

Here we used the CT10 PDFs [40] and the Mathematica package ManeParse [41] to obtain the numerical values of f_q and $f_{\bar{q}}$ at $Q^2 = 20 \text{ GeV}^2$. After including the contributions from scalar, pseudoscalar, and tensor GNI, R^ν and $R^{\bar{\nu}}$ are

$$R^\nu = \frac{(g_L^2 + \frac{1}{3}g_R^2)f_q + (\frac{1}{3}g_L^2 + g_R^2)f_{\bar{q}}}{f_q + \frac{1}{3}f_{\bar{q}}} + \frac{\frac{1}{24} \sum_{q=u,d} ((\epsilon_{s,q}^{\alpha\mu 11})^2 + (\epsilon_{p,q}^{\alpha\mu 11})^2 + 224(\epsilon_{T,q}^{\alpha\mu 11})^2)(f_q + f_{\bar{q}})}{f_q + \frac{1}{3}f_{\bar{q}}}, \quad (3.44)$$

$$R^{\bar{\nu}} = \frac{(\frac{1}{3}g_L^2 + g_R^2)f_q + (g_L^2 + \frac{1}{3}g_R^2)f_{\bar{q}}}{\frac{1}{3}f_q + f_{\bar{q}}} + \frac{\frac{1}{24} \sum_{q=u,d} ((\epsilon_{s,q}^{\alpha\mu 11})^2 + (\epsilon_{p,q}^{\alpha\mu 11})^2 + 224(\epsilon_{T,q}^{\alpha\mu 11})^2)(f_q + f_{\bar{q}})}{\frac{1}{3}f_q + f_{\bar{q}}}. \quad (3.45)$$

Using the numerical values in eq. (3.31) and (3.43), we obtain our naive SM values $R_{SM}^\nu = 0.32$ and $R_{SM}^{\bar{\nu}} = 0.37$, which deviate significantly from the NuTeV measured values in eq. (3.41). Including nuclear effects, partonic charge symmetry violation and strange quarks resolves the NuTeV anomaly [42], bringing the experimental measurements in good agreement with the SM values $R^\nu = 0.3950$ and $R^{\bar{\nu}} = 0.4066$. We simply rescale our naive SM calculations to the more accurate ones. We apply the same rescaling to the new physics contributions to set the 90% C.L. bounds,

$$|\epsilon_{S,q}^{\alpha\mu 11}|, |\epsilon_{P,q}^{\alpha\mu 11}| < 0.19, \quad |\epsilon_{T,q}^{\alpha\mu 11}| < 0.013. \quad (3.46)$$

The degeneracies between the three SMNEFT WCs can be broken by the R^ν and $R^{\bar{\nu}}$ measurements. By plugging the numerical relations in eq. (2.19) into eqs. (3.44) and (3.45), the bounds on the three SMNEFT WCs, allowing all of them to be nonzero simultaneously, are

$$|C_{NLQu}^{\alpha\mu 11}| < 0.078, \quad |C_{NLdQ}^{\alpha\mu 11}| < 0.076, \quad |C_{NLdQ}^{\prime\alpha\mu 11}| < 0.015. \quad (3.47)$$

3.5 CE ν NS

Coherent elastic neutrino-nucleus scattering occurs when the momentum exchanged is smaller than the inverse of the nucleus size, which typically requires neutrino energies of $\mathcal{O}(10 \text{ MeV})$. The cross section is enhanced by the square of the number of nucleons, thus providing an excellent tool to investigate GNI at low energies. The COHERENT experiment has recently observed CE ν NS in a low-threshold CsI detector at the 6.7σ level. This is consistent with the SM at 1σ [43]. The neutrino flux from the Spallation Neutron Source (SNS) is comprised of prompt, monoenergetic ν_μ from stopped pion decays, $\pi^+ \rightarrow \mu^+ + \nu_\mu$, and $\bar{\nu}_\mu$ and ν_e from the subsequent muon decays, $\mu^+ \rightarrow e^+ + \bar{\nu}_\mu + \nu_e$.

The neutrino fluxes are

$$\begin{aligned}\phi_{\nu_\mu}(E_{\nu_\mu}) &= \mathcal{N} \frac{2m_\pi}{m_\pi^2 - m_\mu^2} \delta\left(1 - \frac{2E_{\nu_\mu}m_\pi}{m_\pi^2 - m_\mu^2}\right), \\ \phi_{\nu_e}(E_{\nu_e}) &= \mathcal{N} \frac{192}{m_\mu} \left(\frac{E_{\nu_e}}{m_\mu}\right)^2 \left(\frac{1}{2} - \frac{E_{\nu_e}}{m_\mu}\right), \\ \phi_{\bar{\nu}_\mu}(E_{\bar{\nu}_\mu}) &= \mathcal{N} \frac{64}{m_\mu} \left(\frac{E_{\bar{\nu}_\mu}}{m_\mu}\right)^2 \left(\frac{3}{4} - \frac{E_{\bar{\nu}_\mu}}{m_\mu}\right),\end{aligned}\tag{3.48}$$

where \mathcal{N} is a normalization factor determined by the experimental setup. The ν_μ energy is fixed at $(m_\pi^2 - m_\mu^2)/(2m_\pi) \approx 30 \text{ MeV}$ due to the two-body pion decay. The ν_e and $\bar{\nu}_\mu$ energies have a kinematic upper bound, $m_\mu/2 \approx 50 \text{ MeV}$.

The differential cross section including scalar, vector, and tensor contributions reads [15]

$$\begin{aligned}\frac{d\sigma_a^\beta}{dE_r} &= \frac{G_F^2}{4\pi} M_a N_a^2 \\ &\times \left[(\xi_S^\beta)^2 \frac{E_r}{E_{r,\max}} + (\xi_V^\beta)^2 \left(1 - \frac{E_r}{E_{r,\max}} - \frac{E_r}{E_\nu}\right) + (\xi_T^\beta)^2 \left(1 - \frac{E_r}{2E_{r,\max}} - \frac{E_r}{E_\nu}\right) \right] F^2(q^2),\end{aligned}\tag{3.49}$$

where a denotes the target material and α denotes the neutrino flavor. M_a and N_a are the molar mass of the target nucleus and neutron number of the target, respectively. The flavor index $\beta = \mu$ includes both ν_μ and $\bar{\nu}_\mu$. $F(q^2)$ is the nuclear form factor [44]. The maximum recoil energy $E_{r,\max} = \frac{2E_\nu^2}{M_a + 2E_\nu} \approx \frac{2E_\nu^2}{M_a}$. Since the typical recoil energy E_r is $\mathcal{O}(10) \text{ keV}$, and the neutrino energy E_ν is $\mathcal{O}(10) \text{ MeV}$, we can safely ignore the interference term between scalar and tensor interactions, which is proportional to E_r/E_ν . The ξ_S , ξ_V , and ξ_T collect the contributions from scalar, vector, and tensor interactions,

respectively, and are defined as

$$\begin{aligned}
 (\xi_S^\beta)^2 &= \frac{1}{N_a^2} \left\{ \left(\sum_{q=u,d} 2\text{Re}(\epsilon_{S,q}^{\alpha\beta 11}) \left[N \frac{m_n}{m_q} f_{Tq}^n + Z \frac{m_p}{m_q} f_{Tq}^p \right] \right)^2 \right. \\
 &\quad \left. + \left(\sum_{q=u,d} 2\text{Im}(\epsilon_{S,q}^{\alpha\beta 11}) \left[N \frac{m_n}{m_q} f_{Tq}^n + Z \frac{m_p}{m_q} f_{Tq}^p \right] \right)^2 \right\}, \\
 (\xi_V^\beta)^2 &= \frac{2}{N_a^2} (Z(2g_{V,u} + g_{V,d}) + N(g_{V,u} + 2g_{V,d}))^2, \\
 (\xi_T^\beta)^2 &= \frac{8}{N_a^2} \left(\sum_{q=u,d} 4 \text{Re}(\epsilon_{T,q}^{\alpha\beta 11}) [Z\delta_q^p + N\delta_q^n] \right)^2,
 \end{aligned} \tag{3.50}$$

where f_{Tq}^p and f_{Tq}^n are the mass fractions of quark q in the respective nucleon, and the δ_q 's are the corresponding nucleon tensor charges. The effective vector coupling $g_{V,q}$ is

$$g_{V,q} \equiv g_{L,q} + g_{R,q}. \tag{3.51}$$

The expected number of events per day with recoil energy in the energy range $[E_r, E_r + \Delta E_r]$ and arrival time in the time interval $[t, t + \Delta t]$ is given by

$$N_{th}(t, E_r, \epsilon) = \sum_{\beta=e,\mu} \frac{m_{\text{det}} N_A}{M_a} \int_{\Delta E_r} dE_r \int_{\Delta t} dt \rho_\alpha(t) \int_{E_\nu^{\min}}^{E_\nu^{\max}} dE_\nu \phi_\beta(E_\nu) \frac{d\sigma_a^\beta(\epsilon)}{dE_r}, \tag{3.52}$$

where m_{det} is the detector mass, $N_A = 6.022 \times 10^{23} \text{ mol}^{-1}$, and $\rho_\alpha(t)$ is the arrival time probability density function. To calculate the differential neutrino-nucleus scattering cross section, we need to evaluate the matrix elements of the operators between nuclear states. We adopt the following numerical values of the nuclear matrix elements [45, 46]

$$\begin{aligned}
 f_{Tu}^p &= 0.0208, & f_{Td}^p &= 0.0411, & f_{Tu}^n &= 0.0189, & f_{Td}^n &= 0.0451, \\
 \delta_u^p &= 0.792, & \delta_d^p &= -0.194, & \delta_u^n &= -0.194, & \delta_d^n &= 0.792.
 \end{aligned} \tag{3.53}$$

Following ref. [47], we study the current and projected constraints on the three GNI from the COHERENT experiment. Several COHERENT experiments with multiple targets have been proposed. In this study, we consider a future 750 kg liquid argon (LAr) detector with a 610 kg fiducial mass taking data for four years. The energy threshold is around 20 keV, which is higher than the 6.5 keV CsI energy threshold. The observed event distributions based on the SM simulations are shown in figure 2. The future LAr experiment will provide much more statistics even though it has a higher threshold of nuclear recoil energy. GNI can modify the shape of the recoil energy and temporal distributions. The scalar and tensor GNI distributions comparing to the SM are shown in figure 3. The muon flavor contributions dominate over the electron flavor as there are twice as many muon flavor neutrinos as electron flavor neutrinos. Since the $\bar{\nu}_\mu$ energy distribution peaks at the end point $m_\mu/2$, there are more events in the tail of the energy spectrum for GNI

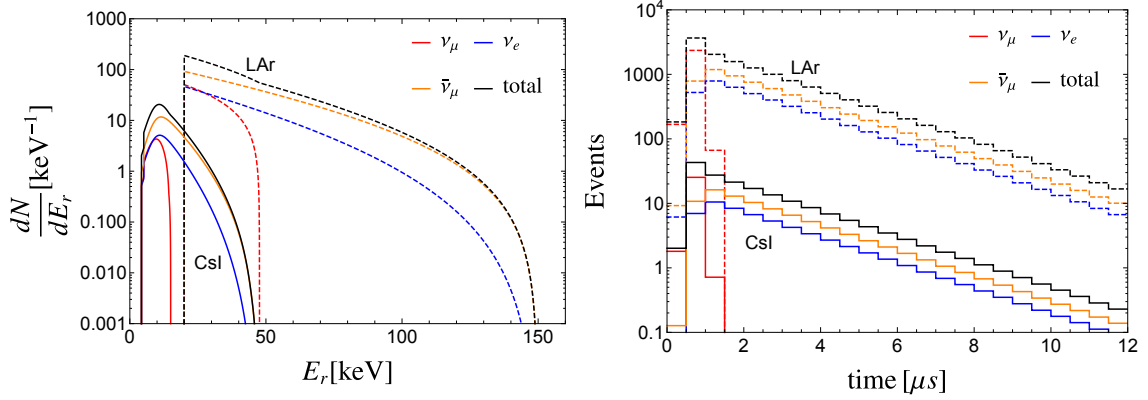


Figure 2. The SM recoil energy (left) and temporal (right) distributions in the current COHERENT CsI detector (solid lines) and a future COHERENT LAr detector (dashed curves). Threshold effects are included. The red (blue) [orange] curves correspond to the contribution from muon (electron) [anti-muon] neutrinos. The black lines correspond to the sum of all the flavor contributions.

involving the muon flavor. Another observation from figure 3 is that COHERENT experiment is much more sensitive to the scalar interactions than tensor interactions. By using the energy spectrum of the current COHERENT data, we find that the current 90% C.L. bounds on the scalar or tensor interactions, allowing only a single nonzero parameter, are

$$(\xi_S^\mu)^2 < 0.60, \quad (\xi_T^\mu)^2 < 0.73, \quad (\xi_S^e)^2 < 1.5, \quad (\xi_T^e)^2 < 1.6. \quad (3.54)$$

Also, the projected 90% C.L. bounds from future COHERENT data by using both the spectral and temporal information are

$$(\xi_S^\mu)^2 < 0.012, \quad (\xi_T^\mu)^2 < 0.013, \quad (\xi_S^e)^2 < 0.030, \quad (\xi_T^e)^2 < 0.027, \quad (3.55)$$

which is an order of magnitude improvement. Again, the bounds are set based on only one of them being nonzero. The projected 90% C.L. bounds in the $(\xi_S^\alpha)^2$ - $(\xi_T^\alpha)^2$ plane are shown in figure 4. Because of degeneracies between the SMNEFT WCs in eq. (2.19), bounds on individual parameters cannot be placed if all the parameters are allowed to float. The bounds on the individual can be derived after running and matching. The current (projected) 90% C.L. bounds on SMNEFT WCs, after setting the others to zero, are

$$\begin{aligned} |C_{NLQu}^{\alpha e 11}| &< 8.1 \times 10^{-2} (3.2 \times 10^{-3}), & |C_{NLQu}^{\alpha \mu 11}| &< 5.1 \times 10^{-2} (2.0 \times 10^{-3}), \\ |C_{NLdQ}^{\alpha e 11}| &< 7.7 \times 10^{-2} (3.1 \times 10^{-3}), & |C_{NLdQ}^{\alpha \mu 11}| &< 4.9 \times 10^{-2} (1.9 \times 10^{-3}), \\ |C_{NLdQ}^{\prime \alpha e 11}| &< 2.0 \times 10^{-1} (2.1 \times 10^{-2}), & |C_{NLdQ}^{\prime \alpha \mu 11}| &< 1.4 \times 10^{-1} (1.4 \times 10^{-2}). \end{aligned} \quad (3.56)$$

The projected 90% C.L. bounds in the C_{NLQu} - C_{NLdQ} (C_{NLdQ} - C_{NLdQ}') planes, are shown by the brown dashed contours in the upper (lower) panels of figure 1. We have set $C_{NLQu} = 0$ in the C_{NLdQ} - C_{NLdQ}' planes, because otherwise the bounds are too weak to display. The current COHERENT bounds are not shown as they are irrelevant in comparison.

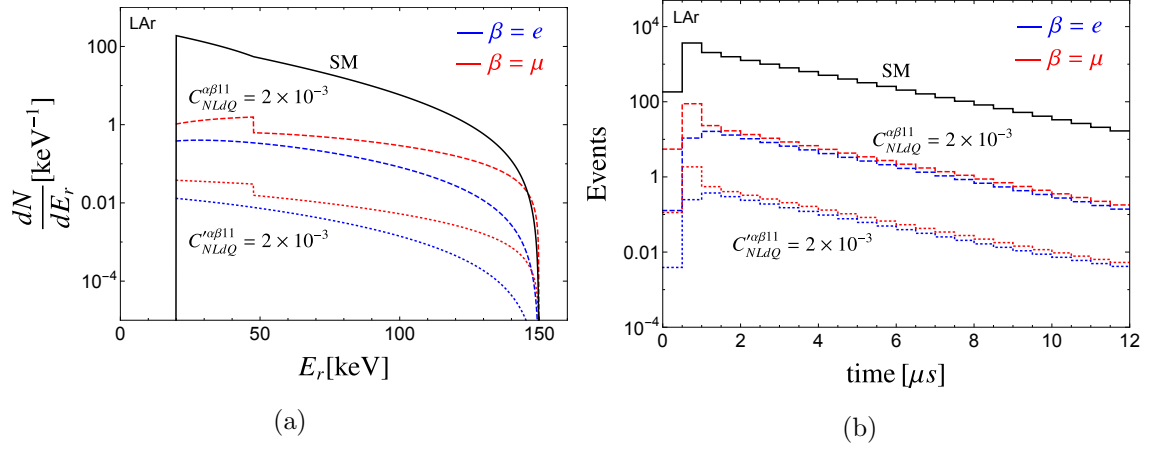


Figure 3. The recoil energy (left) and temporal (right) distributions in a future COHERENT LAr detector. Threshold effects are included. The black solid lines are the SM case including all flavors. The blue (red) curves correspond to the electron (muon+antimuon) flavor contributions. The dashed (dotted) curves correspond to the contributions from the scalar (tensor) interactions with C_{NLdQ} (C'_{NLdQ}) $= 2 \times 10^{-3}$.

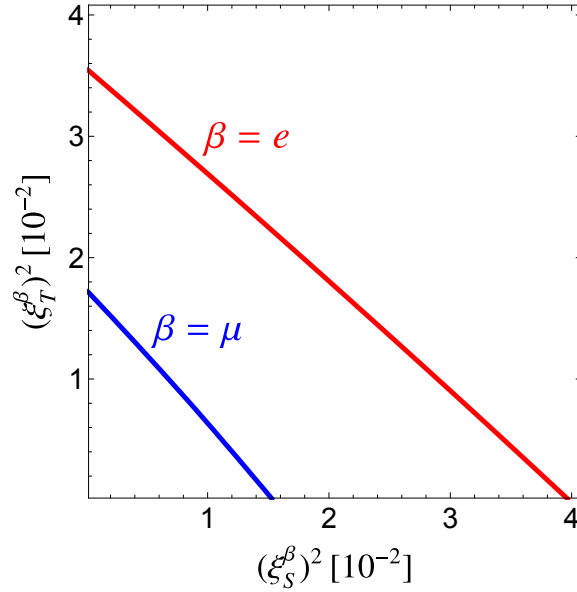


Figure 4. Projected 90% C.L. upper bounds from the future COHERENT experiment with a 610 kg fiducial mass of LAr.

4 Collider constraints

High-energy colliders can set strong bounds on the Wilson coefficients of scalar, pseudoscalar, and tensor interactions. In this section, we study the sensitivity to the WCs at proton-proton and electron-proton colliders. We set bounds using the LHC and evaluate the potential of the HL-LHC and LHeC to probe GNI. By integrating over the full phase space, we find the partonic cross sections of the SM mediated by the W boson and of the contact scalar and tensor interactions to be [48]

$$\text{LHC: } \hat{\sigma}_S = \frac{G_F^2 \hat{s}}{24\pi} C_S^2, \quad \hat{\sigma}_T = \frac{2G_F^2 \hat{s}}{9\pi} C_T^2, \quad \hat{\sigma}_{SM}(u\bar{d} \rightarrow W^* \rightarrow \mu^+ \nu_\mu) = \frac{G_F^2 \hat{s}}{18\pi} \frac{M_W^4}{(\hat{s} - M_W^2)^2}, \quad (4.1)$$

$$\text{LHeC: } \hat{\sigma}_S = \frac{G_F^2 \hat{s}}{24\pi} C_S^2, \quad \hat{\sigma}_T = \frac{14G_F^2 \hat{s}}{3\pi} C_T^2, \quad \hat{\sigma}_{SM}(eq \rightarrow \nu_e q') = \frac{G_F^2 \hat{s}}{2\pi} \frac{M_W^2}{\hat{s} + M_W^2}, \quad (4.2)$$

where $C_S \in \{C_{NLQ}, C_{NLdQ}\}$, $C_T = C'_{NLdQ}$. Note that the ratios of the tensor to scalar cross sections, $\hat{\sigma}_T/\hat{\sigma}_S$, are 16/3 at the LHC and 112 at the LHeC. Clearly, the LHeC is much more sensitive to tensor interactions than scalar interactions. Owing to its lower center-of-mass energy, we expect bounds derived from the LHeC to be weaker than those from the LHC, given the \hat{s} -dependence of the higher-dimensional operators.

The interference between chirality-flipped operators and SM operators are helicity suppressed, and the interference between the scalar (O_{NLdQ}) and tensor (O'_{NLdQ}) interactions is generally nonzero. The differential distributions for the interference of the latter operators in the center-of-mass frame are found to be

$$\begin{aligned} \text{LHC: } \quad \frac{d\hat{\sigma}_{ST}}{d\cos\theta^*} &= \frac{G_F^2 \hat{s}}{12\pi} (C_S^* C_T + C_T^* C_S) \cos\theta^*, \\ \text{LHeC: } \quad \frac{d\hat{\sigma}_{ST}}{d\cos\theta^*} &= \frac{G_F^2 \hat{s}}{16\pi} (C_S^* C_T + C_T^* C_S) (\cos^2\theta^* - 2\cos\theta^* - 3). \end{aligned} \quad (4.3)$$

The interference leads to a linear asymmetry at the LHC and the integrated rate vanishes, while the integrated rate at the LHeC is $\hat{\sigma}_{ST} = -G_F^2 \hat{s}/3\pi (C_S^* C_T + C_T^* C_S)$.

The hadronic cross sections can be obtained by convolving with the parton distribution functions,

$$\sigma_{LHC} = \sum_{q,q'} \int_{\tau_{\min}}^{\tau_{\max}} d\tau \int_{\tau}^1 \frac{dx}{x} f_q(x, \mu_F) f_{q'}(\tau/x, \mu_F) \hat{\sigma}(\tau s), \quad (4.4)$$

$$\sigma_{LHeC} = \sum_q \int_{x_{\min}}^{x_{\max}} dx f_q(x, \mu_F) \hat{\sigma}(xs). \quad (4.5)$$

In the following, we use the Monte Carlo event generator MadGraph5_aMC@NLO [49] to generate signal and background samples at the LHC and LHeC. The GNI Lagrangian is implemented in the FeynRules 2.0 [50] framework. PYTHIA8 [51] (PYTHIA6 [52]) is used for parton showering and hadronization at the LHC (LHeC). We perform the detector simulations using Delphes 3.4.1 [53].

Before evaluating the collider sensitivity to the Wilson coefficients, we note that our EFT description is valid only for $\sqrt{\hat{s}} < \Lambda$, which calls for an assumption about the energy

scale of the new physics. We consider two representative scenarios of the new physics scale, which we call low-scale new physics (LNP) with $\Lambda \sim 1$ TeV, and high-scale new physics (HNP) with $\Lambda \gg 1$ TeV. In the HNP case, we assume the EFT method to be valid for the entire energy scale relevant to LHC data. In the LNP case, however, we limit our analysis to a subset of the LHC data below 1 TeV.

4.1 Proton-proton colliders

Both scalar and tensor CC contact interactions can be probed at high-energy proton-proton colliders, under the assumption that the energy scale of the new dynamics is not kinematically accessible. The signal channel is the Drell-Yan (DY) process, $pp \rightarrow \ell\nu + X$. Due to the missing neutrino in the final state, our analysis is based on the distribution of the transverse mass, which is reconstructed by the charged lepton transverse momentum (p_T^ℓ) and the missing transverse momentum (E_T^{miss})

$$m_T = \sqrt{2p_T^\ell E_T^{\text{miss}}(1 - \cos \Delta\phi(p_T^\ell, E_T^{\text{miss}}))}. \quad (4.6)$$

The main background for large values of m_T is DY production of W bosons. The latest analysis for charged lepton and missing transverse momentum events conducted by ATLAS used 139 fb^{-1} of data collected at $\sqrt{s} = 13$ TeV [54]. In the rest of our study, we only use the m_T distributions below 800 GeV for the LNP scenario, and the full range of m_T for the HNP scenario. For our analyses, we define the statistical significance in terms of

$$\chi^2 = \sum_i \frac{(n_{b,i} + n_{s,i} - n_{\text{data},i})^2}{n_{\text{data},i} + (\sigma_i n_{\text{data},i})^2}, \quad (4.7)$$

where $n_{b(\text{data}),i}$ is the number of background (observed) events in the i^{th} bin, which is obtained directly from ref. [54]. $n_{s,i}$ is the number of signal events simulated in Madgraph at LO. σ_i is the total systematic uncertainty, which is chosen according to ref. [54]:

$$\begin{aligned} \text{electron channel: } \sigma_e &\sim 10\% (12\%) \text{ for } m_T = 300 (2000) \text{ GeV;} \\ \text{muon channel: } \sigma_\mu &\sim 10\% (17\%) \text{ for } m_T = 300 (2000) \text{ GeV.} \end{aligned}$$

The current 90% C.L. bounds, defined by $\Delta\chi^2 < 2.71$, on the LNP (HNP) scalar and tensor operators are

$$\begin{aligned} |C_{NLQ_u}^{\alpha e 11}| &< 2.5 (0.44) \times 10^{-3}, & |C_{NLdQ}^{\alpha e 11}| &< 2.6 (0.46) \times 10^{-3}, \\ |C_{NLdQ}'^{\alpha e 11}| &< 1.2 (0.24) \times 10^{-3}, & & \end{aligned} \quad (4.8)$$

$$\begin{aligned} |C_{NLQ_u}^{\alpha \mu 11}| &< 2.9 (0.66) \times 10^{-3}, & |C_{NLdQ}^{\alpha \mu 11}| &< 3.0 (0.68) \times 10^{-3}, \\ |C_{NLdQ}'^{\alpha \mu 11}| &< 1.4 (0.40) \times 10^{-3}. & & \end{aligned} \quad (4.9)$$

The bounds on C_{NLQ_u} are slightly stronger than for C_{NLdQ} because of the size of the CKM matrix element V_{ud} . These bounds are consistent with those in ref. [54]. The 90% C.L. allowed regions in the C_{NLdQ} - C_{NLQ_u} and C_{NLdQ} - C_{NLdQ}' planes are shown in the figure 1.

The solid red (blue) contours correspond to the LNP (HNP) case. We have checked numerically using Madgraph that the interference between scalar operator O_{NLdQ} and tensor operator O'_{NLdQ} can be ignored.

To assess the future potential of the LHC, we assume an integrated luminosity of $L = 3 \text{ ab}^{-1}$ and $\sqrt{s} = 14 \text{ TeV}$ at the HL-LHC. In this analysis, we simulate the DY W background at LO multiplied by a scale factor obtained from ref. [54], to include other smaller backgrounds including top pairs, single top, $W \rightarrow \tau\nu$, DY Z , and di-bosons. The signals are also generated at tree level. We do not include a K factor as it applies to both signal and background, so the significance is simply scaled by \sqrt{K} after including higher-order corrections. The selection rules applied in this analysis are slightly different between the electron and muon final states. For the muon (electron) final states, we require

- $p_T^{\mu(e)} > 55 (65) \text{ GeV}$ and $|\eta_\ell| < 2.4$,
 - veto b -tagged jets,
 - discard additional electron or muon with $p_T > 20 \text{ GeV}$ and $|\eta_\ell| < 2.4$,
 - $m_T > 300 \text{ GeV}$,
- (4.10)

in which, the electron p_T cut is slightly stronger than the muon p_T cut, in order to suppress the non-prompt backgrounds. The distributions of m_T above 300 GeV after applying the cuts are shown in figure 5a. Deviations from the SM arise in the tails of the m_T distributions because the sub-process cross sections for a dim-6 operator scale as \hat{s} ; see eq. (4.1). For the same size WC, tensor interactions have a larger cross section than scalar interactions. The χ^2 used in this analysis is defined in eq. (4.7), with n_{data} replaced by the values from SM simulations. The projected 90% C.L. bounds on the LNP (HNP) scalar and tensor operators are

$$\begin{aligned}
 |C_{NLQu}^{\alpha e 11}| &< 2.3 (0.28) \times 10^{-3}, & |C_{NLdQ}^{\alpha e 11}| &< 2.4 (0.28) \times 10^{-3}, \\
 |C_{NLdQ}^{\prime \alpha e 11}| &< 1.1 (0.18) \times 10^{-3}, & &
 \end{aligned}
 \tag{4.11}$$

$$\begin{aligned}
 |C_{NLQu}^{\alpha \mu 11}| &< 2.7 (0.28) \times 10^{-3}, & |C_{NLdQ}^{\alpha \mu 11}| &< 2.8 (0.29) \times 10^{-3}, \\
 |C_{NLdQ}^{\prime \alpha \mu 11}| &< 1.3 (0.18) \times 10^{-3}. & &
 \end{aligned}
 \tag{4.12}$$

The bounds from HL-LHC on scalar (tensor) interactions with the assumption of LNP are comparable with (much stronger than) the ones we obtained for the future COHERENT experiment. The dashed red (blue) contours in figure 1 show the 90% C.L. projections for the HL-LHC with the LNP (HNP) assumption. The bounds on the WCs are stronger for HNP than for LNP, because the signals in the high-energy tails of the m_T distributions are not buried in the SM background. These bounds can be converted into limits on the effective couplings $\kappa = \sqrt{|C|}(\Lambda/v)$ for fixed values of the new physics scale Λ . The 90% C.L. bounds on κ are provided in table 2 for LNP (with $\Lambda = 1 \text{ TeV}$) and HNP (with $\Lambda = 10 \text{ TeV}$). As expected, bounds on κ are stronger in the LNP case than the HNP case. Alternatively, if we assume that $\kappa \approx 1$, then HL-LHC bounds on the WCs for HNP imply a sensitivity to $\Lambda \sim 20 \text{ TeV}$. This is comparable to the expected sensitivity of W' searches at the HL-LHC [55].

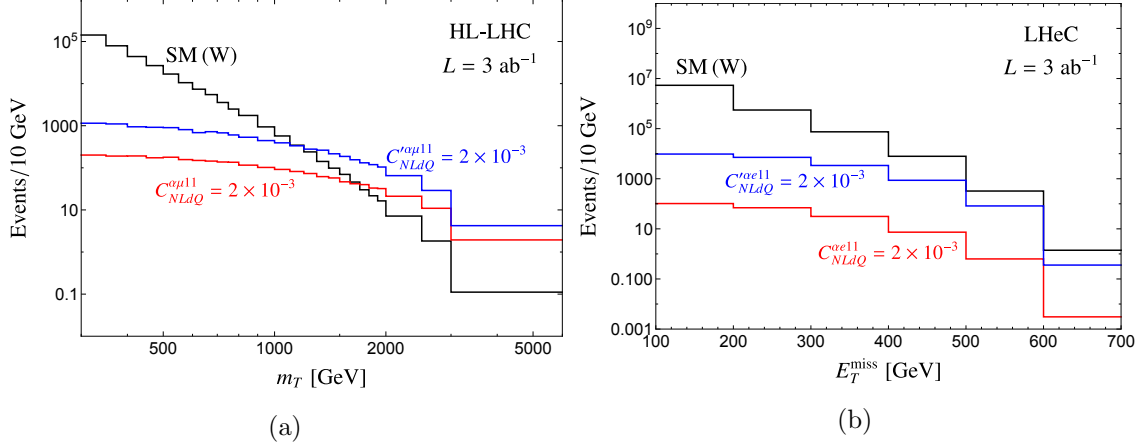


Figure 5. Left: distribution of m_T at the HL-LHC with an integrated luminosity of 3 ab^{-1} . Right: distribution of E_T^{miss} at the LHeC with 3 ab^{-1} and a 1.3 TeV center-of-mass energy. The black histograms corresponds to the SM. The red (blue) histograms correspond to scalar (tensor) interactions with C_{NLdQ} (C'_{NLdQ}) $= 2 \times 10^{-3}$.

Coupling	$\kappa_{NLQu}^{\alpha e11}$	$\kappa_{NLdQ}^{\alpha e11}$	$\kappa'_{NLdQ}^{\alpha e11}$	$\kappa_{NLQu}^{\alpha \mu11}$	$\kappa_{NLdQ}^{\alpha \mu11}$	$\kappa'_{NLQu}^{\alpha \mu11}$
LHC: LNP (HNP)	0.20 (0.85)	0.21 (0.87)	0.14 (0.63)	0.22 (1.0)	0.22 (1.1)	0.15 (0.81)
HL-LHC: LNP (HNP)	0.19 (0.68)	0.20 (0.68)	0.13 (0.55)	0.21 (0.68)	0.22 (0.69)	0.15 (0.55)

Table 2. Current and projected 90% C.L. bounds on the new physics coupling κ from LHC and HL-LHC data, respectively, for the LNP ($\Lambda = 1 \text{ TeV}$) and HNP ($\Lambda = 10 \text{ TeV}$) cases.

4.2 Electron-proton colliders

The HERA collaboration set bounds on the contact interaction $e\nu qq'$ using the charged current process, $e^\pm p \rightarrow (\bar{\nu})X$, from the Q^2 and x distributions [56]. The lower bound on the mass scale of the contact term is around 1 TeV with the strong coupling $\sim 4\pi$. This bound can be translated to our scenario:

$$|C_{NLQu}^{\alpha e11}|, |C_{NLdQ}^{\alpha e11}|, |C'_{NLdQ}^{\alpha e11}| \lesssim 5, \quad (4.13)$$

which is very weak compared to bounds from high-energy colliders.

Next, we consider the future ep collider, LHeC, with $\sqrt{s} = 1.3 \text{ TeV}$ ($E_e = 60 \text{ GeV}$, $E_p = 7 \text{ TeV}$) and $L = 3 \text{ ab}^{-1}$. The signal channel is mono-jet, $ep \rightarrow j\nu + X$, through the t -channel. The main background is mediated by SM W bosons. For the analysis we use the following set of basic cuts:

- leading jet should have $p_T^j > 20 \text{ GeV}$ and $|\eta_j| < 2.5$,
- veto any electrons with $p_T^e > 20 \text{ GeV}$ and $|\eta_e| < 2.5$,
- the angular distance between jet and missing E_T should be bigger than 0.4.

The distributions of the missing transverse energy above 100 GeV after applying the cuts are shown in figure 5b. To maximize our $\chi^2 = S^2/B$, in which we do not include systematic

uncertainties, we select the cut on the missing transverse energy as $E_T^{\text{miss}} > 300 \text{ GeV}$. The projected 90% C.L. bounds on the individual SMNEFT WCs are

$$|C_{NLQu}^{\alpha e 11}| < 3.9 \times 10^{-3}, \quad |C_{NLdQ}^{\alpha e 11}| < 4.0 \times 10^{-3}, \quad |C_{NLdQ}'^{\alpha e 11}| < 0.38 \times 10^{-3}, \quad (4.14)$$

with only one WC taken to be nonzero. If all parameters are allowed to be nonzero, the bounds weaken slightly due to the mixing between O_{NLdQ} and O_{NLdQ}' :

$$|C_{NLQu}^{\alpha e 11}| < 3.9 \times 10^{-3}, \quad |C_{NLdQ}^{\alpha e 11}| < 6.1 \times 10^{-3}, \quad |C_{NLdQ}'^{\alpha e 11}| < 0.58 \times 10^{-3}. \quad (4.15)$$

The projected 90% C.L. bounds on the C_{NLdQ} - C_{NLQu} and C_{NLdQ} - C_{NLdQ}' are shown in figure 1 by the purple dashed contours. Due to the smaller center-of-mass energy, the bounds on the scalar interactions from LHeC are weaker than the ones from HL-LHC. However, for tensor interactions, the bounds from LHeC are stronger than HL-LHC for the LNP case.

5 Summary

New physics associated with neutrinos can be studied without theoretical prejudice, with allowance for scalar, pseudoscalar, vector, axial-vector and tensor interactions of neutrinos with SM fermions. If the new physics scale is much higher than the electroweak scale, it is appropriate to work in a model-independent EFT framework below the new physics scale. GNI operators below the electroweak scale are generated by EFT operators that respect the SM gauge symmetry.

In this work we studied scalar, pseudoscalar and tensor neutrino interactions in the framework of SMNEFT, which extends SMEFT with right-handed neutrinos. At the dimension-six level, these interactions are produced by three less constrained and phenomenologically interesting operators, namely O_{NLQu} , O_{NLdQ} , and O_{NLdQ}' as in section 2.2. Both neutral current and charged current interactions can be induced by a single operator, which can be explored in various experiments. To compare constraints from experiments at different energy scales, we perform the RG running above and below the weak scale, and map all the bounds into the parameter space of three WCs C_{NLQu} , C_{NLdQ} , and C_{NLdQ}' at 1 TeV. We summarize the current and projected experimental bounds on the three WCs in tables 3 and 4. The correlations between the three operators are shown in figure 1. Our main conclusions are:

1. Neutrino mass bounds indicate that the SMNEFT operators involving the second and third families of quarks are highly constrained, while the parameter space for neutrino interactions with the first quark generation is relatively unconstrained. This conclusion, however, is model-dependent and can be evaded.
2. Bounds on the SMNEFT WCs from low-energy probes generally suffer from degeneracies, which are induced by RG running and matching, as is evident from eq. (2.19). The high-energy probes set bounds directly on the SMNEFT WCs, and so are not subject to degeneracies. Low-energy probes and high-energy colliders are complementary.

WC	π^+ decay	β decay	ν DIS	CE ν NS	HERA	LHC: LNP(HNP)
$C_{NLQu}^{\alpha e 11}$	3.3×10^{-6}	3.4×10^{-2}	0.77	8.1×10^{-2}	~ 5	$2.5(0.44) \times 10^{-3}$
$C_{NLdQ}^{\alpha e 11}$	3.4×10^{-6}	3.5×10^{-2}	0.75	7.7×10^{-2}	~ 5	$2.6(0.46) \times 10^{-3}$
$C_{NLdQ}'^{\alpha e 11}$	3.9×10^{-5}	2.8×10^{-2}	0.15	0.20	~ 5	$1.2(0.24) \times 10^{-3}$
$C_{NLQu}^{\alpha \mu 11}$	1.5×10^{-3}	—	7.8×10^{-2}	5.1×10^{-2}	—	$2.9(0.66) \times 10^{-3}$
$C_{NLdQ}^{\alpha \mu 11}$	1.5×10^{-3}	—	7.6×10^{-2}	4.9×10^{-2}	—	$3.0(0.68) \times 10^{-3}$
$C_{NLdQ}'^{\alpha \mu 11}$	1.7×10^{-2}	—	1.5×10^{-2}	0.14	—	$1.4(0.40) \times 10^{-3}$

Table 3. Current 90% C.L. bounds on the three SMNEFT WCs C_{NLQu} , C_{NLdQ} , and C_{NLdQ}' , for the electron and muon flavors at a 1 TeV energy scale. The constraints obtained by allowing all WCs to simultaneously vary are in boldface.

3. Charged pion decay is extremely sensitive to the LEFT pseudoscalar operators. But, there are degeneracies when the bounds are mapped into the SMNEFT WCs. With the assumption of only one nonzero operator at a time, the bounds on the electron flavor are at the 10^{-6} level.
4. The strongest current bounds on the three SMNEFT operators are from LHC charged lepton $+E_T^{\text{miss}}$ searches, and are at the $10^{-4} - 10^{-3}$ level depending on the energy range of validity of the EFT.
5. HL-LHC can improve the bounds by a factor of a few and reach 10^{-4} in the HNP case. For LNP, the improvement is minor because systematic uncertainties dominate for low m_T .
6. Current LHC data can exclude $\kappa \gtrsim 0.14$ for $\Lambda = 1$ TeV and $\kappa \gtrsim 0.63$ for $\Lambda = 10$ TeV. Future HL-LHC data can exclude $\kappa \gtrsim 0.13$ for $\Lambda = 1$ TeV and $\kappa \gtrsim 0.55$ for $\Lambda = 10$ TeV. For strong interactions with $\kappa = 4\pi$, the new physics scale can be excluded up to 200 TeV.
7. A future COHERENT experiment with LAr can set strong bounds on the scalar operators, comparable with that from the HL-LHC with the LNP assumption, especially when the muon flavor is involved.
8. LHeC will be important to study tensor interactions involving the electron flavor, and can place bounds at the 10^{-4} level.

Acknowledgments

We thank Keping Xie for helpful discussions. The work of T.H. and H.L. was supported in part by the U.S. Department of Energy under grant No. DE-FG02-95ER40896 and in part by the PITT PACC. The work of J.L. was supported by the National Natural Science Foundation of China under Grant No. 11905299. D.M. was supported in part by the U.S. Department of Energy under Grant No. de-sc0010504.

WC	CE ν NS-LAr	LHeC	HL-LHC: LNP(HNP)
$C_{NLQu}^{\alpha e 11}$	3.2×10^{-3}	3.9×10^{-3}	$2.3 (0.28) \times 10^{-3}$
$C_{NLdQ}^{\alpha e 11}$	3.1×10^{-3}	6.1×10^{-3}	$2.4 (0.28) \times 10^{-3}$
$C'_{NLdQ}^{\alpha e 11}$	2.1×10^{-2}	0.58×10^{-3}	$1.1 (0.18) \times 10^{-3}$
$C_{NLQu}^{\alpha \mu 11}$	2.0×10^{-3}	—	$2.7 (0.28) \times 10^{-3}$
$C_{NLdQ}^{\alpha \mu 11}$	1.9×10^{-3}	—	$2.8 (0.29) \times 10^{-3}$
$C'_{NLdQ}^{\alpha \mu 11}$	1.4×10^{-2}	—	$1.3 (0.18) \times 10^{-3}$

Table 4. Projected 90% C.L. bounds on the three SMNEFT WCs C_{NLQu} , C_{NLdQ} , and C'_{NLdQ} , with electron and muon flavor, at 1 TeV energy scale. The constraints obtained by allowing all WCs to simultaneously vary are in boldface.

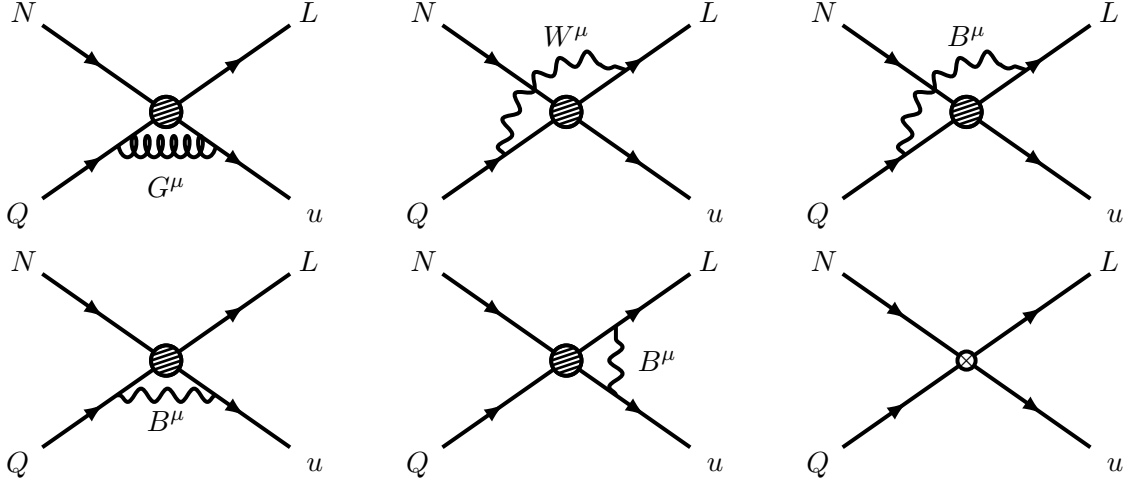


Figure 6. One-loop vertex corrections and counterterm for O_{NLQu} above the weak scale. The diagrams are similar for the other two operators.

A 1-loop RG running

We define the SMNEFT Lagrangian as

$$\mathcal{L}_{\text{SMNEFT}} \supset \sum_{i=1}^3 C_i O_i, \quad (\text{A.1})$$

where $\vec{O} \equiv \{O_{NLQu}, O_{NLdQ}, O'_{NLdQ}\}$ and $\vec{C} \equiv 2\sqrt{2}G_F\{C_{NLQu}, C_{NLdQ}, C'_{NLdQ}\}$. The anomalous dimension of O_i can be obtained by the operator renormalization,

$$\mathcal{L}_{\text{SMNEFT}} \supset \sum_{i=1}^3 C_i \frac{Z_i}{Z_{4\psi,i}} O_i^{(0)}, \quad (\text{A.2})$$

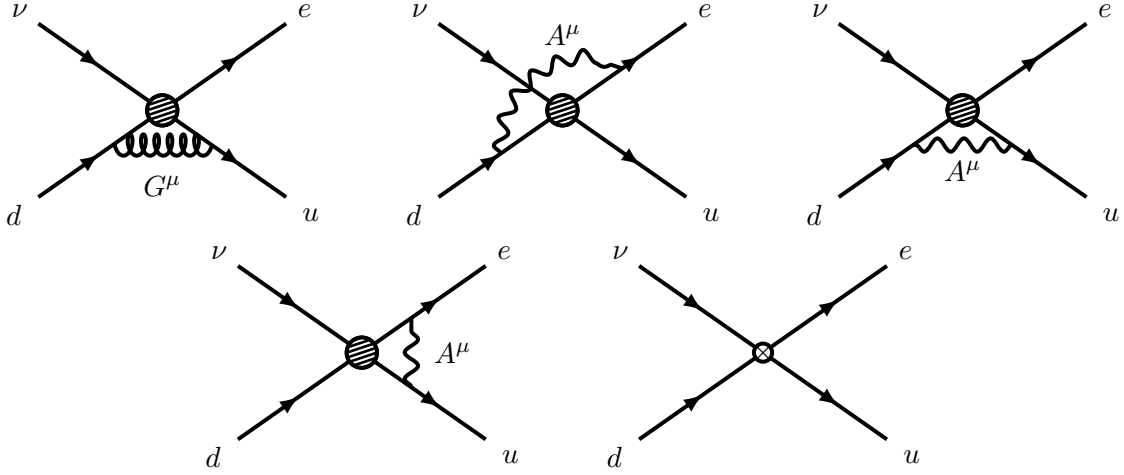


Figure 7. One-loop vertex corrections and counterterm for CC operators.

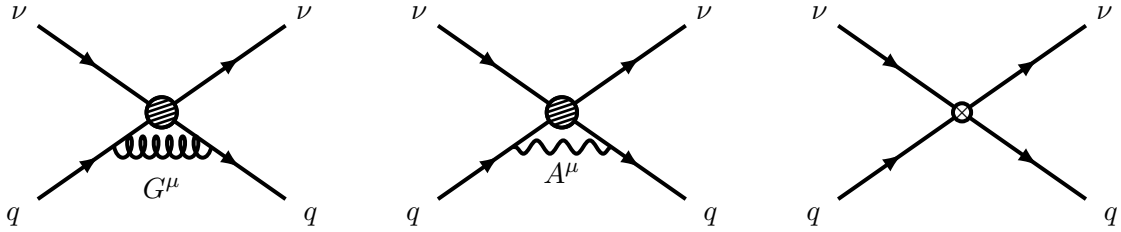


Figure 8. One-loop vertex corrections and counterterm for NC operators.

where $O^{(0)}$ is the operator in terms of bare fields. $Z_{4\psi,i}$ are the four-fermion field strength renormalizations obtained from the self-energy diagrams:

$$Z_{4\psi,1} = \sqrt{Z_{2L}^L Z_{2L}^Y} \sqrt{Z_{2Q}^C Z_{2Q}^L Z_{2Q}^Y} \sqrt{Z_{2u}^C Z_{2u}^Y}, \quad (\text{A.3})$$

$$Z_{4\psi,2} = Z_{4\psi,3} = \sqrt{Z_{2L}^L Z_{2L}^Y} \sqrt{Z_{2Q}^C Z_{2Q}^L Z_{2Q}^Y} \sqrt{Z_{2d}^C Z_{2d}^Y}, \quad (\text{A.4})$$

where the fermion field strengths are

$$Z_{2\psi}^C = 1 - \frac{1}{\epsilon} \frac{2\alpha_3}{3\pi}, \quad Z_{2\psi}^L = 1 - \frac{1}{\epsilon} \frac{3\alpha_2}{8\pi}, \quad Z_{2\psi}^Y = 1 - Y_\psi^2 \frac{1}{\epsilon} \frac{\alpha_1}{2\pi}. \quad (\text{A.5})$$

Here $\alpha_i \equiv \frac{g_i^2}{4\pi}$, and Y_ψ is hypercharge of the fermion. Z_i are the corrections from counterterms which cancels the UV divergence from the 1-loop Feynman diagrams shown in figure 6. We have used FeynCalc [57, 58] to extract the UV divergence from the one-loop integrals. The expressions for Z_i are

$$Z_1 = 1 - \frac{1}{\epsilon} \left[\frac{8\alpha_3}{3\pi} + \frac{3\alpha_2}{8\pi} + \frac{\alpha_1}{2\pi} (Y_u(4Y_Q - Y_L) + Y_L Y_Q) \right], \quad (\text{A.6})$$

$$Z_2 = 1 - \frac{1}{\epsilon} \left[\frac{8\alpha_3}{3\pi} + \frac{3\alpha_2}{8\pi} \left(1 - 12 \frac{C_3}{C_2} \right) + \frac{\alpha_1}{2\pi} \left(Y_L(Y_d - Y_Q) + 4Y_d Y_Q - 12 \frac{C_3}{C_2} \right) \right], \quad (\text{A.7})$$

$$Z_3 = 1 - \frac{1}{\epsilon} \left[\frac{9\alpha_2}{8\pi} \left(1 - \frac{1}{12} \frac{C_2}{C_3} \right) + \frac{3\alpha_1}{2\pi} \left(Y_L(Y_d - Y_Q) + \frac{1}{12} \frac{C_2}{C_3} Y_L(Y_d + Y_Q) \right) \right]. \quad (\text{A.8})$$

Z_2 and Z_3 can introduce mixing between C_2 and C_3 from the weak running. The renormalization group equations of the WCs then arise from the fact that the bare operators and Lagrangian are independent of the renormalization scale μ ,

$$\mu \frac{dC_i}{d\mu} = \gamma_{ij} C_j. \quad (\text{A.9})$$

It is then straightforward to calculate the anomalous dimension matrix γ_{ij} which yield the results shown in eq. (2.11). Following a similar procedure, the anomalous dimension matrix for the CC and NC WCs below the weak scale are obtained by evaluating the one-loop vertex corrections and counterterms in figures 7 and 8, respectively.

Open Access. This article is distributed under the terms of the Creative Commons Attribution License ([CC-BY 4.0](https://creativecommons.org/licenses/by/4.0/)), which permits any use, distribution and reproduction in any medium, provided the original author(s) and source are credited.

References

- [1] S. Weinberg, *A model of leptons*, *Phys. Rev. Lett.* **19** (1967) 1264 [[INSPIRE](#)].
- [2] S.L. Glashow, *Partial symmetries of weak interactions*, *Nucl. Phys.* **22** (1961) 579 [[INSPIRE](#)].
- [3] A. Salam, *Weak and electromagnetic interactions*, *Conf. Proc. C* **680519** (1968) 367 [[INSPIRE](#)].
- [4] V. Barger, D. Marfatia and K. Whisnant, *Progress in the physics of massive neutrinos*, *Int. J. Mod. Phys. E* **12** (2003) 569 [[hep-ph/0308123](#)] [[INSPIRE](#)].
- [5] M.C. Gonzalez-Garcia and M. Maltoni, *Phenomenology with massive neutrinos*, *Phys. Rept.* **460** (2008) 1 [[arXiv:0704.1800](#)] [[INSPIRE](#)].
- [6] L. Wolfenstein, *Neutrino oscillations in matter*, *Phys. Rev. D* **17** (1978) 2369 [[INSPIRE](#)].
- [7] T. Ohlsson, *Status of non-standard neutrino interactions*, *Rept. Prog. Phys.* **76** (2013) 044201 [[arXiv:1209.2710](#)] [[INSPIRE](#)].
- [8] O.G. Miranda and H. Nunokawa, *Non standard neutrino interactions: current status and future prospects*, *New J. Phys.* **17** (2015) 095002 [[arXiv:1505.06254](#)] [[INSPIRE](#)].
- [9] Y. Farzan and M. Tortola, *Neutrino oscillations and non-standard interactions*, *Front. Phys.* **6** (2018) 10 [[arXiv:1710.09360](#)] [[INSPIRE](#)].
- [10] B. Grzadkowski, M. Iskrzynski, M. Misiak and J. Rosiek, *Dimension-six terms in the Standard Model Lagrangian*, *JHEP* **10** (2010) 085 [[arXiv:1008.4884](#)] [[INSPIRE](#)].
- [11] B. Henning, X. Lu and H. Murayama, *How to use the Standard Model effective field theory*, *JHEP* **01** (2016) 023 [[arXiv:1412.1837](#)] [[INSPIRE](#)].
- [12] I. Brivio and M. Trott, *The Standard Model as an effective field theory*, *Phys. Rept.* **793** (2019) 1 [[arXiv:1706.08945](#)] [[INSPIRE](#)].
- [13] Y. Liao and X.-D. Ma, *Operators up to dimension seven in Standard Model effective field theory extended with sterile neutrinos*, *Phys. Rev. D* **96** (2017) 015012 [[arXiv:1612.04527](#)] [[INSPIRE](#)].
- [14] I. Bischer and W. Rodejohann, *General neutrino interactions from an effective field theory perspective*, *Nucl. Phys. B* **947** (2019) 114746 [[arXiv:1905.08699](#)] [[INSPIRE](#)].

- [15] M. Lindner, W. Rodejohann and X.-J. Xu, *Coherent neutrino-nucleus scattering and new neutrino interactions*, *JHEP* **03** (2017) 097 [[arXiv:1612.04150](#)] [[INSPIRE](#)].
- [16] J.C. Pati and A. Salam, *Lepton number as the fourth color*, *Phys. Rev. D* **10** (1974) 275 [*Erratum ibid.* **11** (1975) 703] [[INSPIRE](#)].
- [17] R.N. Mohapatra and J.C. Pati, *A natural left-right symmetry*, *Phys. Rev. D* **11** (1975) 2558 [[INSPIRE](#)].
- [18] G. Senjanović and R.N. Mohapatra, *Exact left-right symmetry and spontaneous violation of parity*, *Phys. Rev. D* **12** (1975) 1502 [[INSPIRE](#)].
- [19] J.C. Pati and A. Salam, *Unified lepton-hadron symmetry and a gauge theory of the basic interactions*, *Phys. Rev. D* **8** (1973) 1240 [[INSPIRE](#)].
- [20] J. Terol-Calvo, M. Tórtola and A. Vicente, *High-energy constraints from low-energy neutrino nonstandard interactions*, *Phys. Rev. D* **101** (2020) 095010 [[arXiv:1912.09131](#)] [[INSPIRE](#)].
- [21] PARTICLE DATA GROUP collaboration, *Review of particle physics*, *Phys. Rev. D* **98** (2018) 030001 [[INSPIRE](#)].
- [22] G. Prezeau and A. Kurylov, *Neutrino mass constraints on μ -decay and $\pi_0 \rightarrow \nu\bar{\nu}$* , *Phys. Rev. Lett.* **95** (2005) 101802 [[hep-ph/0409193](#)] [[INSPIRE](#)].
- [23] T.M. Ito and G. Prezeau, *Neutrino mass constraints on β decay*, *Phys. Rev. Lett.* **94** (2005) 161802 [[hep-ph/0410254](#)] [[INSPIRE](#)].
- [24] A. Loureiro et al., *On the upper bound of neutrino masses from combined cosmological observations and particle physics experiments*, *Phys. Rev. Lett.* **123** (2019) 081301 [[arXiv:1811.02578](#)] [[INSPIRE](#)].
- [25] KATRIN collaboration, *Improved upper limit on the neutrino mass from a direct kinematic method by KATRIN*, *Phys. Rev. Lett.* **123** (2019) 221802 [[arXiv:1909.06048](#)] [[INSPIRE](#)].
- [26] X.-J. Xu, *Tensor and scalar interactions of neutrinos may lead to observable neutrino magnetic moments*, *Phys. Rev. D* **99** (2019) 075003 [[arXiv:1901.00482](#)] [[INSPIRE](#)].
- [27] V. Cirigliano and I. Rosell, *Two-loop effective theory analysis of $\pi(K) \rightarrow e\bar{\nu}_e[\gamma]$ branching ratios*, *Phys. Rev. Lett.* **99** (2007) 231801 [[arXiv:0707.3439](#)] [[INSPIRE](#)].
- [28] V. Cirigliano, S. Gardner and B. Holstein, *Beta decays and non-standard interactions in the LHC era*, *Prog. Part. Nucl. Phys.* **71** (2013) 93 [[arXiv:1303.6953](#)] [[INSPIRE](#)].
- [29] D.I. Britton et al., *Measurement of the $\pi^+ \rightarrow e^+$ neutrino branching ratio*, *Phys. Rev. Lett.* **68** (1992) 3000 [[INSPIRE](#)].
- [30] D.I. Britton et al., *Measurement of the $\pi^+ \rightarrow e^+$ neutrino branching ratio*, *Phys. Rev. D* **49** (1994) 28 [[INSPIRE](#)].
- [31] G. Czapek et al., *Branching ratio for the rare pion decay into positron and neutrino*, *Phys. Rev. Lett.* **70** (1993) 17 [[INSPIRE](#)].
- [32] V. Cirigliano and I. Rosell, *$\pi/K \rightarrow e\bar{\nu}_e$ branching ratios to $O(e^2p^4)$ in chiral perturbation theory*, *JHEP* **10** (2007) 005 [[arXiv:0707.4464](#)] [[INSPIRE](#)].
- [33] B. Ananthanarayan and B. Moussallam, *Four-point correlator constraints on electromagnetic chiral parameters and resonance effective Lagrangians*, *JHEP* **06** (2004) 047 [[hep-ph/0405206](#)] [[INSPIRE](#)].

- [34] S. Descotes-Genon and B. Moussallam, *Radiative corrections in weak semi-leptonic processes at low energy: a two-step matching determination*, *Eur. Phys. J. C* **42** (2005) 403 [[hep-ph/0505077](#)] [[INSPIRE](#)].
- [35] FLAVOUR LATTICE AVERAGING GROUP collaboration, *FLAG review 2019: Flavour Lattice Averaging Group (FLAG)*, *Eur. Phys. J. C* **80** (2020) 113 [[arXiv:1902.08191](#)] [[INSPIRE](#)].
- [36] M. Gonzalez-Alonso, O. Naviliat-Cuncic and N. Severijns, *New physics searches in nuclear and neutron β decay*, *Prog. Part. Nucl. Phys.* **104** (2019) 165 [[arXiv:1803.08732](#)] [[INSPIRE](#)].
- [37] J. Erler and S. Su, *The weak neutral current*, *Prog. Part. Nucl. Phys.* **71** (2013) 119 [[arXiv:1303.5522](#)] [[INSPIRE](#)].
- [38] CHARM collaboration, *Experimental verification of the universality of ν_e and ν_μ coupling to the neutral weak current*, *Phys. Lett. B* **180** (1986) 303 [[INSPIRE](#)].
- [39] NuTeV collaboration, *A precise determination of electroweak parameters in neutrino nucleon scattering*, *Phys. Rev. Lett.* **88** (2002) 091802 [Erratum *ibid.* **90** (2003) 239902] [[hep-ex/0110059](#)] [[INSPIRE](#)].
- [40] H.-L. Lai et al., *New parton distributions for collider physics*, *Phys. Rev. D* **82** (2010) 074024 [[arXiv:1007.2241](#)] [[INSPIRE](#)].
- [41] D.B. Clark, E. Godat and F.I. Olness, *ManeParse: a Mathematica reader for parton distribution functions*, *Comput. Phys. Commun.* **216** (2017) 126 [[arXiv:1605.08012](#)] [[INSPIRE](#)].
- [42] W. Bentz, I.C. Cloet, J.T. Londergan and A.W. Thomas, *Reassessment of the NuTeV determination of the weak mixing angle*, *Phys. Lett. B* **693** (2010) 462 [[arXiv:0908.3198](#)] [[INSPIRE](#)].
- [43] COHERENT collaboration, *Observation of coherent elastic neutrino-nucleus scattering*, *Science* **357** (2017) 1123 [[arXiv:1708.01294](#)] [[INSPIRE](#)].
- [44] S.R. Klein and J. Nystrand, *Interference in exclusive vector meson production in heavy ion collisions*, *Phys. Rev. Lett.* **84** (2000) 2330 [[hep-ph/9909237](#)] [[INSPIRE](#)].
- [45] M. Hoferichter, J. Ruiz de Elvira, B. Kubis and U.-G. Meißner, *High-precision determination of the pion-nucleon σ term from Roy-Steiner equations*, *Phys. Rev. Lett.* **115** (2015) 092301 [[arXiv:1506.04142](#)] [[INSPIRE](#)].
- [46] T. Bhattacharya, V. Cirigliano, S. Cohen, R. Gupta, H.-W. Lin and B. Yoon, *Axial, scalar and tensor charges of the nucleon from $2+1+1$ -flavor lattice QCD*, *Phys. Rev. D* **94** (2016) 054508 [[arXiv:1606.07049](#)] [[INSPIRE](#)].
- [47] T. Han, J. Liao, H. Liu and D. Marfatia, *Nonstandard neutrino interactions at COHERENT, DUNE, T2HK and LHC*, *JHEP* **11** (2019) 028 [[arXiv:1910.03272](#)] [[INSPIRE](#)].
- [48] J. Alcaide, S. Banerjee, M. Chala and A. Titov, *Probes of the Standard Model effective field theory extended with a right-handed neutrino*, *JHEP* **08** (2019) 031 [[arXiv:1905.11375](#)] [[INSPIRE](#)].
- [49] J. Alwall et al., *The automated computation of tree-level and next-to-leading order differential cross sections and their matching to parton shower simulations*, *JHEP* **07** (2014) 079 [[arXiv:1405.0301](#)] [[INSPIRE](#)].

- [50] A. Alloul, N.D. Christensen, C. Degrande, C. Duhr and B. Fuks, *FeynRules 2.0 — a complete toolbox for tree-level phenomenology*, *Comput. Phys. Commun.* **185** (2014) 2250 [[arXiv:1310.1921](#)] [[INSPIRE](#)].
- [51] T. Sjöstrand et al., *An introduction to PYTHIA 8.2*, *Comput. Phys. Commun.* **191** (2015) 159 [[arXiv:1410.3012](#)] [[INSPIRE](#)].
- [52] T. Sjöstrand, S. Mrenna and P.Z. Skands, *A brief introduction to PYTHIA 8.1*, *Comput. Phys. Commun.* **178** (2008) 852 [[arXiv:0710.3820](#)] [[INSPIRE](#)].
- [53] DELPHES 3 collaboration, *DELPHES 3, a modular framework for fast simulation of a generic collider experiment*, *JHEP* **02** (2014) 057 [[arXiv:1307.6346](#)] [[INSPIRE](#)].
- [54] ATLAS collaboration, *Search for a heavy charged boson in events with a charged lepton and missing transverse momentum from pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector*, *Phys. Rev. D* **100** (2019) 052013 [[arXiv:1906.05609](#)] [[INSPIRE](#)].
- [55] ATLAS collaboration, *Prospects for searches for heavy Z' and W' bosons in fermionic final states with the ATLAS experiment at the HL-LHC*, [ATL-PHYS-PUB-2018-044](#), CERN, Geneva, Switzerland (2018).
- [56] F. Cornet and J. Rico, *Contact terms in charged current processes at HERA*, *Phys. Lett. B* **412** (1997) 343 [[hep-ph/9707299](#)] [[INSPIRE](#)].
- [57] V. Shtabovenko, R. Mertig and F. Orellana, *New developments in FeynCalc 9.0*, *Comput. Phys. Commun.* **207** (2016) 432 [[arXiv:1601.01167](#)] [[INSPIRE](#)].
- [58] R. Mertig, M. Böhm and A. Denner, *FEYN CALC: computer algebraic calculation of Feynman amplitudes*, *Comput. Phys. Commun.* **64** (1991) 345 [[INSPIRE](#)].