

Probing neutrino-light-quark effective scalar interactions from neutrino masses

Feng-Zhi Chen[✉], Junlin Huang, and Fanrong Xu^{*}

Department of Physics, College of Physics and Optoelectronic Engineering, Jinan University, Guangzhou 510632, People's Republic of China



(Received 16 December 2025; accepted 11 February 2026; published 6 March 2026)

In this work, we use neutrino masses as a probe of the neutrino-light-quark effective scalar interactions. It is found that neutrinos can acquire masses not only from the usual light quark loop corrections, but also from the light quark condensates. The latter contribution has been overlooked in the literature. We show that both contributions are comparable for operators involving u and d quarks, while quark loop corrections dominate for operators involving the s quark. Using the low-energy effective field theory extended with light right-handed neutrinos and matching to chiral perturbation theory, we systematically analyze these contributions, deriving constraints on the corresponding Wilson coefficients from neutrino mass bounds, coherent elastic neutrino-nucleus scattering, and light pseudoscalar meson invisible decays. Our analysis shows that electron neutrino mass measurements provide the most stringent constraints on these scalar couplings, significantly improving upon limits from other observables. The results highlight the importance of including both perturbative and nonperturbative contributions in complete phenomenological analyses of neutrino mass generation mechanisms.

DOI: [10.1103/gp7s-njbk](https://doi.org/10.1103/gp7s-njbk)

I. INTRODUCTION

In the Standard Model (SM), neutrinos are massless due to the absence of right-handed neutrinos or a Higgs triplet. However, the discovery of neutrino oscillations has demonstrated that neutrinos possess small but nonzero masses [1–3]. Consequently, neutrino masses must originate from new physics (NP) beyond the SM. Investigating the origin of the smallness of neutrino masses provides a promising avenue for exploring the nature of NP.

Experimentally, the most recent terrestrial constraint on the neutrino mass from tritium β decay comes from the KATRIN experiment, which sets an upper limit of $m_{\nu_e} \equiv \sqrt{\sum_i |U_{ei}|^2 m_i^2} < 0.45$ eV at 90% confidence level (CL) [4]. Additionally, stringent bounds on the effective Majorana mass, $m_{ee} < 28\text{--}122$ meV, are derived from the nonobservation of neutrinoless double beta decay by KamLAND-Zen [5]. Cosmological observations of the cosmic microwave background also provide a model-dependent upper bound on the sum of neutrino masses, $\sum_i m_{\nu_i} < 0.12$ eV (95% CL) [6]. In all cases, neutrino masses are at least 6 orders of magnitude smaller than

those of all other SM fermions, suggesting that neutrinos acquire mass through a mechanism distinct from the Higgs mechanism [7,8].

This paper focuses on constraining the effective couplings of neutrino-light-quark scalar interactions using neutrino mass bounds. We assume that neutrinos acquire mass via scalar four-fermion effective interactions involving neutrinos and light quarks (u , d , s). Within this framework, we identify two distinct contributions to neutrino masses: (i) perturbative contributions from light quark loop corrections and (ii) nonperturbative contributions from light quark condensates. While neutrinos could also gain mass from scalar interactions involving heavy quarks or charged leptons via loop diagrams, we focus on light quarks due to their unique property: Below the chiral symmetry-breaking scale ($\mu_\chi \sim 1$ GeV), light quarks not only contribute perturbatively, but also develop a vacuum condensate, providing an additional contribution to the neutrino mass [9–11]. Moreover, the same scalar operators contributing to neutrino masses also contribute to low-energy processes such as coherent elastic neutrino-nucleus scattering (CE ν NS) and invisible decays of light pseudo-scalar mesons.¹ Thus, constraints derived from neutrino masses have significant implications for these processes.

^{*}Contact author: fanrongxu@jnu.edu.cn

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International license](https://creativecommons.org/licenses/by/4.0/). Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

¹While these scalar operators may modify neutrino oscillations (see, e.g., Refs. [12–14]), such effects are expected to be subdominant and are, thus, excluded from our phenomenological analysis.

It is noteworthy that previous studies have considered only the perturbative contribution, overlooking the nonperturbative effect [15–17]. We demonstrate that both contributions are comparable for operators involving u and d quarks, whereas the perturbative contribution dominates for operators involving the s quark. We discuss both Dirac and Majorana neutrino cases, as the nature of neutrinos remains unknown. The analysis is performed using the low-energy effective field theory extended with right-handed neutrinos N_R (LNEFT) [18–21], and we include renormalization group (RG) running effects. The nonperturbative contributions are obtained by matching LNEFT operators onto chiral perturbation theory (χ PT) operators.

The paper is organized as follows. Section II introduces the scalar interactions in the effective field theories, including LNEFT, RG running effects, and the matching to χ PT. Section III presents the expressions for relevant observables: neutrino masses, CE ν NS, and light pseudo-scalar meson invisible decays. Section IV provides a numerical analysis comparing the perturbative and non-perturbative contributions to neutrino masses and constrains the Wilson coefficients from various observables. Our conclusions are given in Sec. V.

II. SCALAR INTERACTIONS IN EFFECTIVE FIELD THEORIES

In this section, we will first describe all the relevant scalar effective interactions in the Lagrangian of LNEFT at scale $\mu_\chi \sim 1$ GeV, from which one can calculate the perturbative contribution to the neutrino masses. We subsequently match the LNEFT operators onto the χ PT operators to obtain the nonperturbative contribution to the neutrino masses.

A. Scalar interactions in LNEFT

LNEFT is an effective theory valid below the electroweak scale (μ_{EW}), invariant under Lorentz transformation and $SU(3)_C \times U(1)_{em}$ gauge symmetries. Its degrees of freedom include all SM fields except the heavy W , Z , h , and t , along with light right-handed neutrinos N_R [18–21]. We consider only the scalar four-fermion operators involving neutrinos and light quarks.

Depending on lepton number conservation, these operators are categorized into two classes: lepton-number-conserving (LNC) and lepton-number-violating (LNV). For Dirac neutrinos, the relevant LNC Lagrangian is

$$\begin{aligned} \mathcal{L}_{\text{LNEFT}}^{\text{LNC}} \supset & L_{\nu Nu}^{S,RR}(\bar{\nu}_{Lp} N_{Rr})(\bar{u}_{Ls} u_{Rt}) + L_{\nu Nu}^{S,RL}(\bar{\nu}_{Lp} N_{Rr})(\bar{u}_{Rs} u_{Lt}) \\ & + L_{\nu Nd}^{S,RR}(\bar{\nu}_{Lp} N_{Rr})(\bar{d}_{Ls} d_{Rt}) + L_{\nu Nd}^{S,RL}(\bar{\nu}_{Lp} N_{Rr})(\bar{d}_{Rs} d_{Lt}) + \text{H.c.}, \end{aligned} \quad (1)$$

where the coupling constant L_i in front of each four-fermion operator is the Wilson coefficient, which is suppressed by a factor of μ^{-2} , L and R denote separately the left and right chirality, and p , r , s , and t denote the generation indices. Since both ν_L and N_R carry a lepton number: $L(\nu_L) = L(N_R) = +1$, interactions governed by Eq. (1) correspond to the LNC ($|\Delta L| = 0$) processes. For Majorana neutrinos, the LNV Lagrangian is

$$\begin{aligned} \mathcal{L}_{\text{LNEFT}}^{\text{LNV}} \supset & L_{\nu\nu}^{S,LR}(\bar{\nu}_{Lp}^c \nu_{Lr})(\bar{u}_{Ls} u_{Rt}) + L_{\nu\nu}^{S,LL}(\bar{\nu}_{Lp}^c \nu_{Lr})(\bar{u}_{Rs} u_{Lt}) \\ & + L_{\nu d}^{S,LR}(\bar{\nu}_{Lp}^c \nu_{Lr})(\bar{d}_{Ls} d_{Rt}) + L_{\nu d}^{S,LL}(\bar{\nu}_{Lp}^c \nu_{Lr})(\bar{d}_{Rs} d_{Lt}) \\ & + L_{\nu u}^{S,RR}(\bar{N}_{Rp}^c N_{Rr})(\bar{u}_{Ls} u_{Rt}) + L_{\nu u}^{S,RL}(\bar{N}_{Rp}^c N_{Rr})(\bar{u}_{Rs} u_{Lt}) \\ & + L_{\nu d}^{S,RR}(\bar{N}_{Rp}^c N_{Rr})(\bar{d}_{Ls} d_{Rt}) + L_{\nu d}^{S,RL}(\bar{N}_{Rp}^c N_{Rr})(\bar{d}_{Rs} d_{Lt}) + \text{H.c.}, \end{aligned} \quad (2)$$

where the charge conjugation of a left-handed field $\psi_L^c = C\bar{\psi}_L^T$ ($C = i\gamma_2\gamma_0$) is actually a right-handed field, and vice versa. Since the antiparticle of a Majorana neutrino is the particle itself, operators in Eq. (2) violate the lepton number by two units ($|\Delta L| = 2$), interactions governed by which correspond to the LNV processes.

B. RG running effects

Since there is a large mass gap between scales μ_χ and μ , one has to resum the large logarithms produced by the ratio of the two scales in the perturbative expansion. This can be

obtained solving the relevant RG equations. Note that LNEFT is merely valid below μ_{EW} , at the range between μ_{EW} and μ ; however, the SM gauge symmetries are still unbroken, and physics at this range should be described by the SM effective field theory [22–25] extended with right-handed neutrinos (SMNEFT) [26–31]. The Lagrangian of SMNEFT is invariant under Lorentz transformation and $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetries, and the degrees of freedom of which are all the SM fields and the extended right-handed neutrinos N_R . To connect the physics at μ_χ to the UV-complete model at μ , a standard procedure is describing the low-energy physics in terms of LNEFT at μ_χ ,

evolving the LNEFT Wilson coefficients from μ_χ up to μ_{EW} , matching LNEFT onto SMNEFT at μ_{EW} , running the SMNEFT Wilson coefficients from μ_{EW} up to μ , and matching SMNEFT to the UV-complete model at μ . Nevertheless, as we are interested in only the neutrino mass generation mechanism and the neutrino low-energy physics, strictly executing such a procedure will be rather tricky and is out of the scope of our research. For simplicity, we neglect the RG evolutions for SMNEFT Wilson coefficients from μ down to μ_{EW} and consider only the one-loop RG running effects from μ_{EW} down to μ_χ for LNEFT Wilson coefficients in this work. This treatment is reasonable, as the dominant running effect is mainly obtained from solving the QCD one-loop RG equations between μ_{EW} and μ_χ , with a result by multiplying the Wilson coefficients at μ_{EW} a factor of about 2. By the way, we also neglect the potential operator mixing that could appear in either LNEFT or SMNEFT, which is equal to assume that the contribution produced only by scalar operators.

Working in the $\overline{MS}$ renormalization scheme, the one-loop QCD and QED RG equations for the scalar LNEFT Wilson coefficients read

$$\mu \frac{d}{d\mu} L_{iq'prqq}^{S,AB} = - \left(6C_F \frac{\alpha_s}{4\pi} + 6Q_q^2 \frac{\alpha}{4\pi} \right) L_{iq'prqq}^{S,AB}, \quad i = \nu, N, \nu N, \quad (3)$$

where $A, B = L, R$, $q' = u, d$ (when $q' = u$, $q = u$, and when $q' = d$, $q = d$ or s), $\alpha_s = g_s^2/4\pi$ and $\alpha = e^2/4\pi$ denote separately the coupling constants of strong and electromagnetic interactions, Q_q is the electric charge for quark with $Q_u = +2/3$ and $Q_d = Q_s = -1/3$, and $C_F = (N_C^2 - 1)/2N_C = 4/3$ with $N_C = 3$ being the color number of QCD. The solution of Eq. (3) is

$$L_{iq'prqq}^{S,AB}(\mu_1) = \left[\frac{\alpha_s(\mu_2)}{\alpha_s(\mu_1)} \right]^{-\frac{6C_F}{2\beta_0}} \left[\frac{\alpha(\mu_2)}{\alpha(\mu_1)} \right]^{-\frac{6Q_q^2}{2b_{0,e}}} L_{iq'prqq}^{S,AB}(\mu_2), \quad (4)$$

where $\beta_0 = 11 - 2n_f/3$ with n_f being the number of active quarks between μ_1 and μ_2 and $b_{0,e} = -4(3n_\ell + 4n_u + n_d)/9$ with n_ℓ , n_u , and n_d being, respectively, the numbers of active charged leptons, up-type quarks, and down-type quarks between μ_1 and μ_2 . Numerically, by quoting the input $\alpha_s(m_Z) = 0.1180(9)$ and $\alpha = 1/137.036$ from PDG [32], one has

$$L_{iq'prqq}^{S,AB}(1 \text{ GeV}) \approx 2.07 L_{iq'prqq}^{S,AB}(m_Z). \quad (5)$$

Clearly, the scalar Wilson coefficients get enhanced when running from a high scale down to a low scale. Note that the running effects of QED are negligible compared to that of QCD, which has also been pointed out in Ref. [21].

C. Matching LNEFT to χ PT

For energy below $\mu_\chi \sim 4\pi F_\pi$ with $F_\pi = 92.3(1)$ MeV being the pion decay constant [33], the observed degrees of freedom are no longer the light quarks but the pseudoscalars (π, K, η), which can be well described by the so-called χ PT. Therefore, to describe physics below μ_χ with contributions from high energies, the LNEFT operators should be matched onto the ones of χ PT.

The χ PT is a QCD low-energy effective field theory, in which the pseudoscalars are the Goldstone bosons corresponding to the spontaneous global chiral symmetries breaking: $SU(3)_L \times SU(3)_R \rightarrow SU(3)_V$. There are produced eight pseudoscalars, which can be written into the theory in a nonlinear way:

$$U = \exp \left(i \frac{\Phi}{F_0} \right), \quad \Phi = \lambda^a \phi^a = \begin{pmatrix} \pi^0 + \frac{1}{\sqrt{3}}\eta & \sqrt{2}\pi^+ & \sqrt{2}K^+ \\ \sqrt{2}\pi^- & -\pi^0 + \frac{1}{\sqrt{3}}\eta & \sqrt{2}K^0 \\ \sqrt{2}K^- & \sqrt{2}\bar{K}^0 & -\frac{2}{\sqrt{3}}\eta \end{pmatrix}, \quad (6)$$

where F_0 is the pion decay constant in the chiral limit, ϕ^a are the Goldstone bosons, and λ^a ($a = 1, \dots, 8$) are the Gell-Mann matrices satisfying the trace relation $\text{Tr}(\lambda_a \lambda_b) = 2\delta_{ab}$. Under $SU(3)_L \times SU(3)_R$ symmetries, the matrix U transforms as $U \rightarrow RUL^\dagger$, with $L(R)$ being the representation of $SU(3)_{L(R)}$. The chiral symmetries are broken by the quark vacuum expectation values or, equivalently, the quark condensates, $\langle 0|\bar{q}q|0\rangle$ ($q = u, d, s$), whose values can be determined through the lattice QCD calculations [34,35]. It is worth pointing out that the quark masses softly break the $SU(3)_L \times SU(3)_R$ symmetries which would result in nonzero masses pseudoscalars.

Beside the pure QCD, one is usually dealing with interactions between quark bilinears and other fields, e.g., the neutrinos along with the accompanying Wilson coefficients in Eqs. (1) and (2). These fields can be treated as the external sources (which are also referred as the spurion fields [36]), which can be organized into the Lagrangian by employing the external field method [37–39]:

$$\begin{aligned} \mathcal{L} &= \mathcal{L}_{\text{QCD}}^0 + \mathcal{L}_{\text{ext}} \\ &= \mathcal{L}_{\text{QCD}}^0 + \bar{q}_L \gamma^\mu l_\mu q_L + \bar{q}_R \gamma^\mu r_\mu q_R + \bar{q}_L S q_R + \bar{q}_R S^\dagger q_L \\ &\quad + \bar{q}_L \sigma^{\mu\nu} t_{\mu\nu} q_R + \bar{q}_R \sigma^{\mu\nu} t_{\mu\nu}^\dagger q_L, \end{aligned} \quad (7)$$

where $q = (u, d, s)^T$ and l_μ , r_μ , S , and $t_{\mu\nu}$ are 3×3 Hermitian matrices in flavor space, which stand for the left-handed, right-handed, scalar, and tensor external sources, respectively. In writing down Eq. (7), we have endowed the external sources with the following chiral power counting:

$$l_\mu \sim \mathcal{O}(p), \quad r_\mu \sim \mathcal{O}(p), \quad S \sim \mathcal{O}(p^2), \quad t_{\mu\nu} \sim \mathcal{O}(p^2). \quad (8)$$

For convenience, the external sources can split into two components [36]:

$$l_\mu \mapsto l_\mu + \tilde{l}_\mu, \quad r_\mu \mapsto r_\mu + \tilde{r}_\mu, \quad S \mapsto S + \tilde{S}, \quad t_{\mu\nu} \mapsto t_{\mu\nu} + \tilde{t}_{\mu\nu}, \quad (9)$$

where S , l_μ , r_μ , and $t_{\mu\nu}$ in the right side of each arrow denote the quark mass matrix and the couplings to electromagnetic field A_μ in the SM:

$$S \mapsto -M^\dagger, \quad l_\mu \mapsto -eQA_\mu, \quad r_\mu \mapsto -eQA_\mu, \quad t_{\mu\nu} \mapsto 0, \quad (10)$$

with $M = \text{diag}(m_u, m_d, m_s)$ and $Q = \text{diag}(2/3, -1/3, -1/3)$; whereas $\tilde{S}$, $\tilde{l}_\mu$, $\tilde{r}_\mu$, and $\tilde{t}_{\mu\nu}$ encode the contributions from NP.

As we are interested in only the scalar interactions between quarks and neutrinos, we merely focus on the term $\bar{q}_L \tilde{S} q_R$

(and its Hermitian conjugate $\bar{q}_R \tilde{S}^\dagger q_L$) in Eq. (7). Concretely, the $\tilde{S}$ matrix in the flavor space can be written as

$$\tilde{S} = \begin{pmatrix} \tilde{S}_{uu} & \tilde{S}_{ud} & \tilde{S}_{us} \\ \tilde{S}_{du} & \tilde{S}_{dd} & \tilde{S}_{ds} \\ \tilde{S}_{su} & \tilde{S}_{sd} & \tilde{S}_{ss} \end{pmatrix}. \quad (11)$$

Therefore, by matching $\bar{q}_L \tilde{S} q_R$ to the LNC LNEFT operators in Eq. (1), one has

$$\tilde{S}_{uu} = L_{\nu Nu}^{S,RR}(\bar{\nu}_{Lp} N_{Rr}) + L_{\nu Nu}^{S,RL*}(\bar{N}_{Rr} \nu_{Lp}), \quad (12)$$

$$\tilde{S}_{dd} = L_{\nu Nd}^{S,RR}(\bar{\nu}_{Lp} N_{Rr}) + L_{\nu Nd}^{S,RL*}(\bar{N}_{Rr} \nu_{Lp}), \quad (13)$$

$$\tilde{S}_{ss} = L_{\nu Ns}^{S,RR}(\bar{\nu}_{Lp} N_{Rr}) + L_{\nu Ns}^{S,RL*}(\bar{N}_{Rr} \nu_{Lp}). \quad (14)$$

Similarly, by matching $\bar{q}_L \tilde{S} q_R$ to the LNV LNEFT operators in Eq. (2), one obtains

$$\tilde{S}_{uu} = L_{\nu Nu}^{S,LR}(\bar{\nu}_{Lp}^c \nu_{Lr}) + L_{\nu Nu}^{S,LL*}(\bar{\nu}_{Lr} \nu_{Lp}^c) + L_{\nu Nu}^{S,RR}(\bar{N}_{Rp}^c N_{Rr}) + L_{\nu Nu}^{S,RL*}(\bar{N}_{Rr} N_{Rp}^c), \quad (15)$$

$$\tilde{S}_{dd} = L_{\nu Nd}^{S,LR}(\bar{\nu}_{Lp}^c \nu_{Lr}) + L_{\nu Nd}^{S,LL*}(\bar{\nu}_{Lr} \nu_{Lp}^c) + L_{\nu Nd}^{S,RR}(\bar{N}_{Rp}^c N_{Rr}) + L_{\nu Nd}^{S,RL*}(\bar{N}_{Rr} N_{Rp}^c), \quad (16)$$

$$\tilde{S}_{ss} = L_{\nu Ns}^{S,LR}(\bar{\nu}_{Lp}^c \nu_{Lr}) + L_{\nu Ns}^{S,LL*}(\bar{\nu}_{Lr} \nu_{Lp}^c) + L_{\nu Ns}^{S,RR}(\bar{N}_{Rp}^c N_{Rr}) + L_{\nu Ns}^{S,RL*}(\bar{N}_{Rr} N_{Rp}^c). \quad (17)$$

On the other hand, the χ PT counterparts of $\bar{q}_L \tilde{S} q_R$ at the lowest chiral power order $[\mathcal{O}(p^4)]$ read [36]

$$\begin{aligned} \bar{q}_L \tilde{S} q_R \rightarrow & -2B_0 \left[\frac{1}{4} F_0^2 \langle \tilde{S} U \rangle + L_4 \langle D_\mu U^\dagger D^\mu U \rangle \langle \tilde{S} U \rangle + L_5 \langle \tilde{S} U D_\mu U^\dagger D^\mu U \rangle \right. \\ & \left. + 2L_6 \langle U^\dagger \chi + \chi^\dagger U \rangle \langle \tilde{S} U \rangle - 2L_7 \langle U^\dagger \chi - \chi^\dagger U \rangle \langle \tilde{S} U \rangle + 2L_8 \langle \tilde{S} U \chi^\dagger U \rangle + H_2 \langle \tilde{S} \chi \rangle \right] + \mathcal{O}(p^6), \end{aligned} \quad (18)$$

where $\langle \cdots \rangle$ represents the trace in flavor space, $\chi = -2B_0 S^\dagger$ with $B_0 = -\langle 0 | \bar{q} q | 0 \rangle / F_0^2$, and L_i ($i = 4, 5, 6, 7, 8$) and H_2 the low-energy coupling constants of the corresponding terms. It is clear from Eq. (18) that this term is proportional to a constant B_0 , which is actually connected to the quark condensates, and is, therefore, a nonperturbative effect. Expanding the first term on the right side of Eq. (18) at the lowest energy, we will obtain an effective mass term for neutrinos; for more details, see Sec. III A. By the way, it is noteworthy that a similar procedure has also been applied to obtain the nonperturbative effects in other observables, e.g., the charged lepton and neutrino electric or magnetic dipole moments [36,40].

III. OBSERVABLES

In this section, we will in detail represent the formulas for the relevant observables, which are expressed in terms

of the LNEFT Wilson coefficients. These expressions will be used for the numerical analysis in the next section.

A. Neutrino masses

As introduced in Sec. I, the neutrinos acquire mass through two distinct ways: (i) the perturbative contribution from light quark loop corrections and (ii) the nonperturbative contribution from light quark condensates. In the following, we will separately display the detailed formulas for these two sources.

The perturbative contribution is produced from the Feynman diagrams depicted in Fig. 1.² These types of Feynman diagrams were first calculated in Refs. [15,16],

²Note that the scalar effective interactions should be produced at a scale above μ_{EW} , so that the heavy degree of freedom Z can appear in the two-loop diagrams.

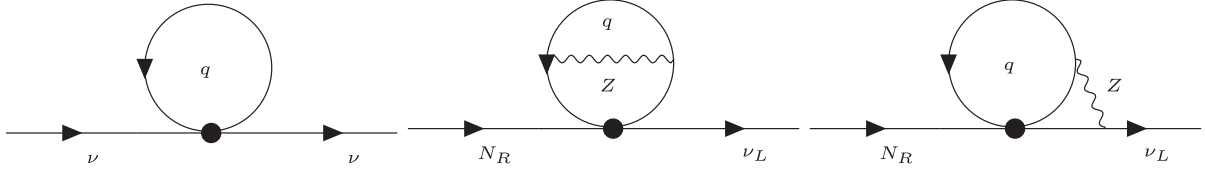


FIG. 1. One- and two-loop Feynman diagrams contributing to the Dirac neutrino masses, where the blob denotes the neutrino-light-quark scalar effective interactions. Similar Feynman diagrams also apply to generate Majorana neutrino masses.

with an effective neutrino mass at one- and two-loop levels given, respectively, by

$$\delta m_{\text{eff}}^{(1)} \approx \frac{1}{2} N_C L_{\text{pr}qq}^{S,AB} \frac{m_q^3}{(4\pi)^2} \log \frac{\mu^2}{m_q^2}, \quad (19)$$

$$\delta m_{\text{eff}}^{(2)} \approx \frac{1}{2} g^2 N_C L_{\text{pr}qq}^{S,AB} \frac{m_q m_Z^2}{(4\pi)^4} \left(\log \frac{\mu^2}{m_Z^2} \right)^2, \quad (20)$$

where g is the $SU(2)_L$ gauge coupling constant. It is found that the ratio of the one-loop result to the two-loop result is $\sim 50 m_q^2/m_Z^2$, so the one-loop contribution is negligible compared to the two-loop one; see more details about this aspect in Refs. [15,16]. Note that, since N_R is an $SU(2)_L \times U(1)_Y$ singlet, only the first two diagrams have contribution to the pure right-handed neutrino mass, yet for the sake of uniformity we keep the same formalism for all the neutrino mass terms. The relations between the Wilson coefficients of this work and the ones in Ref. [15] are

$$\begin{aligned} L_{\nu N q'}^{S,RR} &= 2G_F (a_{S;pr}^{qq} + a_{SP;pr}^{qq}), \\ L_{\nu N q'}^{S,RL} &= 2G_F (a_{S;pr}^{qq} - a_{SP;pr}^{qq}), \end{aligned} \quad (21)$$

$$\begin{aligned} L_{\nu q'}^{S,LR} &= 2G_F (a_{S;pr}^{qq} - a_{SP;pr}^{qq}), \\ L_{\nu q'}^{S,LL} &= 2G_F (a_{S;pr}^{qq} + a_{SP;pr}^{qq}), \end{aligned} \quad (22)$$

where G_F is the Fermi constant and the subscripts S and SP refer to scalar and pseudoscalar, respectively. In the numerical analysis, we consider only the predominant two-loop result, and the total perturbative contribution is obtained by multiplying Eq. (20) with the RG running effect in Eq. (4).

The nonperturbative contribution, on the other hand, can be obtained from the first term in Eq. (18), which corresponds to the contribution from the quark condensate. To see this, we substitute the matching relations in Eqs. (12)–(17) into their χ PT counterparts (18) and take only the vacuum state of the theory, i.e., the 3×3 identity matrix of the U matrix ($U = \mathbb{1} + i\Phi/F_0 + \dots$), to obtain the following effective mass term:

$$\delta m_{\text{eff}}^{qc} = -\frac{1}{2} B_0 F_0^2 \sum_q \tilde{S}_{qq} = \frac{1}{2} \langle 0 | \bar{q}q | 0 \rangle \sum_q L_{\text{pr}qq}^{S,AB}. \quad (23)$$

The numerical value of $\langle 0 | \bar{q}q | 0 \rangle$ will be quoted from the results of lattice QCD [34,35]. The total neutrino effective mass is the sum of both the perturbative and nonperturbative contributions. We will numerically compare these two effects in Sec. IV.

B. CE ν NS

In a CE ν NS process, the low-energy (anti)neutrinos can coherently couple to protons and neutrons of a given nucleus $\mathcal{N}$, the cross section of which could significantly get enhanced and, therefore, allows for much smaller detectors (see, e.g., the recent proceedings on the CE ν NS measurement in the CONUS+ experiment [41]). For this, it has been theoretically proposed long ago to use the CE ν NS to probe the weak neutral current [42]. Experimentally, the CE ν NS process was observed by the COHERENT experiment first in cesium iodide [43] and later confirmed in argon [44] and germanium [45]; all are consistent with the SM predictions. One, therefore, can estimate impact on the CE ν NS with influence from various NP sources [21,46,47]. As the same LNEFT scalar operators contributing to neutrino masses can also contribute to the CE ν NS, we will use the results of the Wilson coefficients constrained from the former to predict their impacts on the latter.

The incoming (anti)neutrinos in a CE ν NS process can be definitely selected, but the outgoing missing energies cannot be measured and might be either the left- or right-handed neutrinos. This allows us to set up a limit on the NP parameters that could have a contribution to the process. In the presence of scalar interactions, the differential cross section for the scattering $\bar{\nu} \mathcal{N} \rightarrow X \mathcal{N}$ with $X \in \{ \bar{\nu}^{(-)}, N^{(-)} \}$ can be written as [21,46–48]

$$\frac{d\sigma}{dT} = \frac{G_F^2 M}{4\pi} \left[\xi_S^2 \frac{E_r}{E_r^{\max}} + \xi_V^2 \left(1 - \frac{E_r}{E_r^{\max}} \right) \right], \quad (24)$$

where M is the mass of the given nucleus, E_r denotes recoil nucleus kinetic energy with maximal value $E_r^{\max} = 2E_\nu^2/(M + E_\nu^2) \simeq 2E_\nu^2/M$ with E_ν being the energy of the incoming (anti)neutrinos, and ξ_S and ξ_V define the effective parameters describing the neutrino-nucleus interactions for scalar and vector currents, respectively. The SM contributes only to the vector parameter ξ_V , with

$$\xi_{V,SM}^2 = [\mathbb{N} - (1 - 4\sin^2\theta_W)\mathbb{Z}]^2 F^2(q^2), \quad (25)$$

where $\mathbb{Z}$ and $\mathbb{N}$ label separately the proton and neutron numbers of nucleus $\mathcal{N}$ and $F(q^2) = F_p(q^2) = F_n(q^2)$

denotes the Helm form factor of the nucleus with the coherent limit $F(q^2 \rightarrow 0) \simeq 1$ [49]. The parameter ξ_S encodes the contributions from the scalar effective interactions, according to different participating neutrinos; one has [21,46,47]

$$\xi_S^2 = \begin{cases} \sum_{j:pr} \left| \frac{1}{\sqrt{2}G_F} \sum_q (L_{iq'}^{S,RR} + L_{iq'}^{S,RL}) \left(\mathbb{Z}_j \frac{m_p}{m_q} f_{T_q}^p + \mathbb{N}_j \frac{m_n}{m_q} f_{T_q}^n \right) \right|^2 F^2(q^2), & \text{for } i = \nu N, \\ \sum_{j:pr} \left| \frac{\sqrt{2}}{G_F} \sum_q (L_{iq'}^{S,LR} + L_{iq'}^{S,LL}) \left(\mathbb{Z}_j \frac{m_p}{m_q} f_{T_q}^p + \mathbb{N}_j \frac{m_n}{m_q} f_{T_q}^n \right) \right|^2 F^2(q^2), & \text{for } i = \nu, \end{cases} \quad (26)$$

where the subscript j sums over the nucleus of a given molecule that participates the scattering, m_p and m_n denote, respectively, the mass of proton and neutron, and $f_{T_q}^{p,n}$ are the nucleon form factors for the scalar currents, with numeric values [50,51]

$$\begin{aligned} f_{T_u}^p &= 0.0153, & f_{T_d}^p &= 0.0191, & f_{T_s}^p &= 0.0447, \\ f_{T_u}^n &= 0.0110, & f_{T_d}^n &= 0.0273, & f_{T_s}^n &= 0.0447. \end{aligned} \quad (27)$$

C. Light pseudoscalar meson invisible decays

In the SM, the electric neutral light pseudoscalar (π^0 , η , and η') which consists of light quarks can decay into a neutrino pair via exchanging a Z gauge boson. The SM predictions for these decay rates are highly suppressed, and only the upper limits were obtained from experiments, which, at 90% CL [52,53],

$$\begin{aligned} \mathcal{B}(\pi^0 \rightarrow \text{invisible}) &< 4.4 \times 10^{-9}, \\ \mathcal{B}(\eta \rightarrow \text{invisible}) &< 1.1 \times 10^{-4}, \\ \mathcal{B}(\eta' \rightarrow \text{invisible}) &< 2.1 \times 10^{-4}, \end{aligned} \quad (28)$$

where the missing energy are supposed to be carried by the invisible neutrinos. As the missing energy may be mimicked by other exotic light degrees of freedom, the aforementioned light pseudoscalar meson invisible decays can be also used to set bounds on the potential NP parameters.

In the presence of scalar interactions, the decay rate of a pseudoscalar meson P (where $P \in \{\pi^0, \eta, \eta'\}$) invisible decay reads [21]

$$\begin{aligned} \mathcal{B}(P \rightarrow \text{invisible}) &= \frac{\tau_P m_P}{16\pi} \sum_{q:pr} \left[2 \left| \frac{h_P^q}{4m_q} (L_{\nu N_R q'}^{S,RR} - L_{\nu N_R q'}^{S,RL}) \right|^2 \left(1 - \frac{m_{N_R}^2}{m_P^2} \right)^2 \right. \\ &\quad \left. + \left| \frac{h_P^q}{2m_q} (L_{N_R q'}^{S,RL} - L_{N_R q'}^{S,RR}) \right|^2 \left(1 - 2 \frac{m_{N_R}^2}{m_P^2} \right) \left(1 - 4 \frac{m_{N_R}^2}{m_P^2} \right)^{\frac{1}{2}} + \left| \frac{h_P^q}{2m_q} (L_{\nu q'}^{S,LL} - L_{\nu q'}^{S,LR}) \right|^2 \right], \end{aligned} \quad (29)$$

where m_{N_R} , m_q , and m_P denote the mass of the right-handed neutrino N_R , constituent quark q , and pseudoscalar meson P , respectively, and h_P^q the decay constant which parametrized the hadronization of P via a pseudoscalar current transition into vacuum:

$$\langle 0 | \bar{q} \gamma_5 q | P \rangle = -i \frac{h_P^q}{2m_q}. \quad (30)$$

Especially,

$$\begin{aligned} h_\pi^u &= m_\pi^2 f_\pi^u, & h_\pi^d &= m_\pi^2 f_\pi^d, & h_\pi^s &= 0, \\ h_\eta^u &= \frac{c_\phi}{\sqrt{2}} h_q, & h_\eta^d &= \frac{c_\phi}{\sqrt{2}} h_q, & h_\eta^s &= -s_\phi h_s, \\ h_{\eta'}^u &= \frac{s_\phi}{\sqrt{2}} h_q, & h_{\eta'}^d &= \frac{s_\phi}{\sqrt{2}} h_q, & h_{\eta'}^s &= c_\phi h_s, \end{aligned} \quad (31)$$

where $s_\phi \equiv \sin \phi$ and $c_\phi \equiv \cos \phi$ with $\phi = 39.3^\circ$ being the $\eta - \eta'$ mixing angle in the Feldmann-Kroll-Stech scheme [54–56] and

$$h_q = f_q(m_\eta^2 c_\phi^2 + m_{\eta'}^2 s_\phi^2) - \sqrt{2} f_s(m_{\eta'}^2 - m_\eta^2) s_\phi c_\phi,$$

$$h_s = f_s(m_\eta^2 c_\phi^2 + m_{\eta'}^2 s_\phi^2) - \frac{f_q}{\sqrt{2}}(m_{\eta'}^2 - m_\eta^2) s_\phi c_\phi \quad (32)$$

with $f_\pi^u = -f_\pi^d = F_\pi$, $f_q = 1.51F_\pi$, and $f_s = 1.90F_\pi$. Note that, in the SM, only vector operators which arise from exchanging a Z gauge boson contribute to the light pseudoscalar meson decays. However, due to that their contributions to the decay are helicity suppressed, such contributions are negligible [21].

IV. NUMERICAL ANALYSIS

In this section, we will first numerically compare the perturbative and nonperturbative contributions to the neutrino mass, then constrain and compare the LNEFT Wilson coefficients from the neutrino masses, $\text{CE}\nu\text{NS}$, and (π^0, η, η') pseudoscalar meson invisible decays, and finally use the results obtained from neutrino mass bounds as inputs to predict the light pseudoscalar meson decay rates. For convenience, the numerical values for all inputs (and the corresponding references) used throughout this work are collected in Table I.

A. Perturbative contribution versus nonperturbative contribution

As demonstrated in Sec. III A, the LNEFT neutrino-light-quark scalar operators have two contributions to the effective neutrino masses: one is the perturbative contribution arising from loop corrections, and the other is the nonperturbative contribution originating from the quark condensates. Using inputs from Table I, the total effective neutrino mass for a single nonzero Wilson coefficient is

$$\delta m_\nu^{\text{eff}} = \begin{cases} (0.022_{\text{loop}} - 0.011_{qc}) L_{pq}^{S,AB} \cdot \text{GeV}^3, & \text{for } q = u, \\ (0.047_{\text{loop}} - 0.011_{qc}) L_{pq}^{S,AB} \cdot \text{GeV}^3, & \text{for } q = d, \\ (0.944_{\text{loop}} - 0.013_{qc}) L_{pq}^{S,AB} \cdot \text{GeV}^3, & \text{for } q = s, \end{cases} \quad (33)$$

where “loop” and “qc” above indicate the contribution induced from loop corrections and quark condensates, respectively. Obviously, we can conclude from Eq. (33) that both contributions are comparable for u and d quarks, while the perturbative contribution dominates for s quarks. Therefore, in the former case, both contributions should be taken into account, while we need to consider only the perturbative contribution in the latter case for which the nonperturbative contribution is negligible. Note that only the perturbative contribution has been considered in Ref. [15], so their results should be refined. Besides, the updated bounds on the neutrino masses also demand renewed constraints on the scalar couplings.

It is important to note that the bounds derived in Eq. (33) rely on a naturalness argument. We require that the contribution to the neutrino mass induced by the neutrino-light-quark scalar effective interactions does not exceed the experimentally observed neutrino mass. Mathematically, it is possible for these effective couplings to be significantly larger if the effective operator contributions are cancelled by a dimension-5 Weinberg operator (or other contributions) with opposite sign. Such a scenario would require a precise fine-tuning between unrelated model parameters. In the absence of a symmetry enforcing such a cancellation, we consider the naturalness bounds to be the most physically robust constraints.

While we have analyzed the neutrino mass generation within the framework of LNEFT, the embedding of these operators into a UV-complete model requires careful consideration of matching effects. In the context of SMNEFT, integrating out heavy scalars or vectors to generate dimension-6 operators typically induces a contribution to the dimension-5 Weinberg operator (or the Dirac Yukawa couplings and Majorana masses in models with right-handed neutrinos) at the one-loop matching level.

From a power-counting perspective, the matching coefficient for an operator of dimension d scales as $\frac{C^{(d)}}{\mu^{d-4}} \sim \frac{y^n}{(16\pi^2)^k \mu^{d-4}}$, where y represents the underlying UV couplings, n denotes the coupling power depending on the specific Feynman diagram, and k stands for the loop order. For a high UV scale μ , the matching contribution to

TABLE I. The numerical values for the relevant inputs used in our numerical analysis.

$G_F = 1.1663788 \times 10^{-5} \text{ GeV}^{-2}$	[32]	$\alpha_s(m_Z) = 0.1180(9)$	[32]
$\alpha = 1/137.035999084(21)$	[32]	$\sin^2 \theta_W = 0.23129(4)$	[32]
$m_{\pi^0} = 134.9768(5) \text{ MeV}$	[32]	$\tau_{\pi^0} = 8.43(13) \times 10^{-17} \text{ s}$	[32]
$m_\eta = 547.862(17) \text{ MeV}$	[32]	$\Gamma_\eta = 1.31(5) \text{ keV}$	[32]
$m_{\eta'} = 957.78(6) \text{ MeV}$	[32]	$\Gamma_{\eta'} = 0.188(6) \text{ MeV}$	[32]
$m_u = 2.16(4) \text{ MeV}$	[32]	$m_d = 4.70(4) \text{ MeV}$	[32]
$m_s = 93.5(5) \text{ MeV}$	[32]	$m_c = 1.2730(28) \text{ GeV}$	[32]
$m_b = 4.183(4) \text{ GeV}$	[32]	$m_Z = 91.1880(20) \text{ GeV}$	[32]
$\langle 0 \bar{u}u 0\rangle^{\overline{\text{MS}}}(2 \text{ GeV}) = -(283(2) \text{ MeV})^3$	[34]	$\langle 0 \bar{d}d 0\rangle^{\overline{\text{MS}}}(2 \text{ GeV}) = -(283(2) \text{ MeV})^3$	[34]
$\langle 0 \bar{s}s 0\rangle^{\overline{\text{MS}}}(2 \text{ GeV}) = -(296(11) \text{ MeV})^3$	[35]	$F_\pi = 92.3(1) \text{ MeV}$	[33]

TABLE II. Constraints on the upper limits of LNC scalar LNEFT Wilson coefficients from $m_{\nu_e}^{\text{eff}}$, $\text{CE}\nu\text{NS}$, and light pseudoscalar meson invisible decays, respectively.

L_i [GeV^{-2}]	$m_{\nu_e}^{\text{eff}}$	$\text{CE}\nu\text{NS}$	$\pi^0 \rightarrow \text{inv}$	$\eta \rightarrow \text{inv}$	$\eta' \rightarrow \text{inv}$
$L_{\nu Nu}^{S,RR(RL)}$	4.3×10^{-8}	7.6×10^{-7}	4.8×10^{-7}	8.5×10^{-4}	1.5×10^{-2}
$L_{\nu Nd}^{S,RR(RL)}$	1.2×10^{-8}	8.8×10^{-7}	1.1×10^{-6}	1.8×10^{-3}	3.2×10^{-2}
$L_{\nu Nd}^{S,RR(RL)}$	4.8×10^{-10}	9.4×10^{-6}	$\dots$	5.5×10^{-4}	5.6×10^{-3}

the neutrino mass from the dimension-5 operator generally dominates over the contribution from the dimension-6 operators, unless the underlying flavor symmetries of the UV theory suppress the dimension-5 operator while allowing the dimension-6 interactions.

B. Constraints on the Wilson coefficients

Since the same LNEFT Wilson coefficients simultaneously contribute to neutrino masses, $\text{CE}\nu\text{NS}$, and pseudoscalar meson invisible decays, different upper limits on these Wilson coefficients can be obtained from various observables. In this section, we will in detail compare the constraint capacity of different observables.

Although our effective neutrino mass formulas apply to neutrinos of all flavors, we merely focus on the electron neutrino ν_e , for which the constraints from the muon- (ν_μ) and tau- (ν_τ) neutrinos masses are still very poor, and the most stringent upper bound to date is from the ν_e mass. As the contributions from the scalar interactions of neutrinos and light quarks are only part of the corrections to the total neutrino masses, we have the following constraint:

$$\delta m_{\nu}^{\text{eff}} = \delta m_{\nu_{\text{eff}}}^{(2)} + \delta m_{\nu_{\text{eff}}}^{qc} < m_{\nu}^{\text{eff}}. \quad (34)$$

For the effective Dirac ν_e mass, we will use the 90% CL upper bound from KATRIN Collaboration [4]:

$$m_{\nu_e}^{\text{eff}} < 0.45 \text{ eV}. \quad (35)$$

While for the effective Majorana electron neutrino mass, the strongest upper limit is taken from the bound on the $0\nu\beta\beta$ decay lifetime for xenon from the KamLAND-Zen experiment [5], which at 90% CL is given by

$$m_{ee}^{\text{eff}} < 28\text{--}122 \text{ meV}. \quad (36)$$

There has been no direct measurement on the light right-handed neutrino masses, so no available neutrino mass data can be utilized to constrain the relevant Wilson coefficients. However, the light pseudoscalar meson invisible decays can provide an indirect constraint on these Wilson coefficients, which, in turn, can be translated into the upper limit of correction to the mass of the right-handed neutrino.

In the following, we will compare the upper limits on the LNEFT Wilson coefficients obtained, respectively, from neutrino masses, $\text{CE}\nu\text{NS}$, and light pseudoscalar meson invisible decays. All the numerical values correspond to the LNEFT Wilson coefficients at the scale $\mu = 1 \text{ GeV}$ and are obtained with the assumption that one coefficient is non-zero at a time. The constraints from the latter two have been done in Ref. [21] in the limit of massless right-handed neutrinos; we will quote their bounds from $\text{CE}\nu\text{NS}$ based on the COHERENT measurement with cesium iodide [43] but use the updated constraints from the light pseudoscalar meson invisible decays. The comparisons of upper limits on LNC and LNV scalar LNEFT Wilson coefficients obtained from various observables are displayed in detail in Tables II and III, respectively. From these two tables, it is clear that neutrino mass bounds provide the most stringent

TABLE III. Constraints on the upper limits of LNV scalar LNEFT Wilson coefficients from m_{ee}^{eff} , $\text{CE}\nu\text{NS}$, and light pseudoscalar meson invisible decays, respectively.

L_i [GeV^{-2}]	m_{ee}^{eff}	$\text{CE}\nu\text{NS}$	$\pi^0 \rightarrow \text{inv}$	$\eta \rightarrow \text{inv}$	$\eta' \rightarrow \text{inv}$
$L_{\nu u}^{S,LL(LR)}$	$(1.3\text{--}5.8) \times 10^{-9}$	3.8×10^{-7}	3.0×10^{-7}	6.0×10^{-4}	1.1×10^{-2}
$L_{\nu d}^{S,LL(LR)}$	$(3.9\text{--}16.9) \times 10^{-10}$	4.4×10^{-7}	6.7×10^{-7}	1.3×10^{-3}	2.3×10^{-2}
$L_{\nu d}^{S,LL(LR)}$	$(1.5\text{--}6.6) \times 10^{-11}$	4.7×10^{-6}	$\dots$	3.9×10^{-4}	4.0×10^{-3}
$L_{\nu u}^{S,RR(RL)}$	$\dots$	$\dots$	3.0×10^{-7}	6.0×10^{-4}	9.2×10^{-3}
$L_{\nu d}^{S,RR(RL)}$	$\dots$	$\dots$	6.7×10^{-7}	1.3×10^{-3}	2.0×10^{-2}
$L_{\nu d}^{S,RR(RL)}$	$\dots$	$\dots$	$\dots$	3.9×10^{-4}	4.0×10^{-3}

constraints on both the LNC and LNV scalar LNEFT Wilson coefficients, surpassing those from CE ν NS and light pseudoscalar meson invisible decays.

From Table III, we see that, among all of the light pseudoscalar meson invisible decays, the π^0 invisible decay provides the most stringent upper bound on the Wilson coefficients $L_{euu}^{S,RR(RL)}$ and $L_{edd}^{S,RR(RL)}$, while the η invisible decay imposes the most strict upper limit on $L_{eess}^{S,RR(RL)}$. As discussed above, these upper limits can be translated into a correction to the mass of the light right-handed neutrino, with

$$\delta m_{N_R}^{\text{eff}} < \begin{cases} 6.6 \text{ eV}, & \text{from } L_{euu}^{S,RR(RL)}, \\ 48.2 \text{ eV}, & \text{from } L_{edd}^{S,RR(RL)}, \\ 0.7 \text{ MeV}, & \text{from } L_{eess}^{S,RR(RL)}. \end{cases} \quad (37)$$

Future more precise measurement on these light pseudoscalar meson invisible decays will improve the current upper bounds on the mass of the light right-handed neutrino.

V. CONCLUSION

This work has systematically investigated the constraints on neutrino-light-quark scalar couplings from neutrino mass measurements, considering both perturbative loop corrections and nonperturbative quark condensate contributions. We have shown that, for operators involving u and d quarks, both contributions are comparable and must be included in complete analyses, while for s quark operators, perturbative effects dominate.

By comparing constraints from various observables, we find that electron neutrino mass measurements provide the most stringent limits on the relevant LNEFT Wilson

coefficients, significantly outperforming constraints from CE ν NS and light pseudoscalar meson invisible decays. For Dirac neutrinos, the effective electron neutrino mass bound of 0.45 eV constrains the couplings to approximately 10^{-8} – 10^{-10} GeV $^{-2}$, while for Majorana neutrinos, the more stringent 28–122 meV bound from neutrinoless double beta decay provides constraints at the 10^{-9} – 10^{-11} GeV $^{-2}$ level. It should be stressed that in deriving these constraints we have assumed the absence of accidental fine-tuned cancellations between the neutrino-light-quark scalar effective operator contributions and other contributions to the neutrino mass matrix, such as the dimension-5 Weinberg operator.

Furthermore, we have translated the constraints from light pseudoscalar meson invisible decays into limits on possible right-handed neutrino masses, obtaining bounds ranging from eV to MeV scales depending on the specific operator. Our results update previous analyses that considered only perturbative contributions and demonstrate the importance of a comprehensive treatment including both perturbative and nonperturbative effects.

The methodology developed in this work provides a framework for connecting UV-complete models with low-energy observables through effective field theory techniques, with potential applications to other beyond-Standard Model scenarios involving neutrino-mass generation mechanisms.

ACKNOWLEDGMENTS

This work is supported in part by NSFC under Grants No. 12475095 and No. U1932104 and the 2024 Guangzhou Basic and Applied Basic Research Scheme Project for Maiden Voyage (2024A04J4190).

DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

-
- [1] Y. Fukuda *et al.* (Super-Kamiokande Collaboration), *Phys. Rev. Lett.* **81**, 1562 (1998).
 - [2] Q. R. Ahmad *et al.* (SNO Collaboration), *Phys. Rev. Lett.* **89**, 011301 (2002).
 - [3] K. Eguchi *et al.* (KamLAND Collaboration), *Phys. Rev. Lett.* **90**, 021802 (2003).
 - [4] M. Aker *et al.* (KATRIN Collaboration), *Science* **388**, adq9592 (2025).
 - [5] S. Abe *et al.* (KamLAND-Zen Collaboration), *Phys. Rev. Lett.* **135**, 262501 (2025).
 - [6] N. Aghanim *et al.* (Planck Collaboration), *Astron. Astrophys.* **641**, A6 (2020); **652**, C4(E) (2021).
 - [7] P. W. Higgs, *Phys. Rev. Lett.* **13**, 508 (1964).
 - [8] F. Englert and R. Brout, *Phys. Rev. Lett.* **13**, 321 (1964).
 - [9] S. D. Thomas and R.-M. Xu, *Phys. Lett. B* **284**, 341 (1992).
 - [10] A. Babič, S. Kovalenko, M. I. Krivoruchenko, and F. Šimkovic, *Phys. Rev. D* **103**, 015007 (2021).
 - [11] O. Scholer, On the implications of a future neutrinoless double beta decay discovery, Ph.D. thesis, U. Heidelberg (main), 2024.
 - [12] S.-F. Ge and S. J. Parke, *Phys. Rev. Lett.* **122**, 211801 (2019).
 - [13] P. B. Denton, A. Giarnetti, and D. Meloni, *J. High Energy Phys.* **02** (2023) 210.
 - [14] P. B. Denton, A. Giarnetti, and D. Meloni, *J. High Energy Phys.* **01** (2025) 097.

- [15] G. Prezeau and A. Kurylov, *Phys. Rev. Lett.* **95**, 101802 (2005).
- [16] T.M. Ito and G. Prezeau, *Phys. Rev. Lett.* **94**, 161802 (2005).
- [17] T. Han, J. Liao, H. Liu, and D. Marfatia, *J. High Energy Phys.* **07** (2020) 207.
- [18] E. E. Jenkins, A. V. Manohar, and P. Stoffer, *J. High Energy Phys.* **03** (2018) 016; **12** (2023) 43.
- [19] E. E. Jenkins, A. V. Manohar, and P. Stoffer, *J. High Energy Phys.* **01** (2018) 084; **12** (2023) 42.
- [20] M. Chala and A. Titov, *J. High Energy Phys.* **05** (2020) 139.
- [21] T. Li, X.-D. Ma, and M. A. Schmidt, *J. High Energy Phys.* **07** (2020) 152.
- [22] W. Buchmuller and D. Wyler, *Nucl. Phys.* **B268**, 621 (1986).
- [23] B. Grzadkowski, M. Iskrzynski, M. Misiak, and J. Rosiek, *J. High Energy Phys.* **10** (2010) 085.
- [24] L. Lehman, *Phys. Rev. D* **90**, 125023 (2014).
- [25] I. Brivio and M. Trott, *Phys. Rep.* **793**, 1 (2019).
- [26] F. del Aguila, S. Bar-Shalom, A. Soni, and J. Wudka, *Phys. Lett. B* **670**, 399 (2009).
- [27] A. Aparici, K. Kim, A. Santamaria, and J. Wudka, *Phys. Rev. D* **80**, 013010 (2009).
- [28] S. Bhattacharya and J. Wudka, *Phys. Rev. D* **94**, 055022 (2016); **95**, 039904(E) (2017).
- [29] Y. Liao and X.-D. Ma, *J. High Energy Phys.* **11** (2016) 043.
- [30] Y. Liao and X.-D. Ma, *Phys. Rev. D* **96**, 015012 (2017).
- [31] I. Bischer and W. Rodejohann, *Nucl. Phys.* **B947**, 114746 (2019).
- [32] S. Navas *et al.* (Particle Data Group), *Phys. Rev. D* **110**, 030001 (2024).
- [33] A. Manohar and H. Georgi, *Nucl. Phys.* **B234**, 189 (1984).
- [34] C. McNeile, A. Bazavov, C. T. H. Davies, R. J. Dowdall, K. Hornbostel, G. P. Lepage, and H. D. Trottier, *Phys. Rev. D* **87**, 034503 (2013).
- [35] C. T. H. Davies, K. Hornbostel, J. Komijani, J. Koponen, G. P. Lepage, A. T. Lytle, and C. McNeile (HPQCD Collaboration), *Phys. Rev. D* **100**, 034506 (2019).
- [36] W. Dekens, E. E. Jenkins, A. V. Manohar, and P. Stoffer, *J. High Energy Phys.* **01** (2019) 088.
- [37] J. Gasser and H. Leutwyler, *Ann. Phys. (N.Y.)* **158**, 142 (1984).
- [38] J. Gasser and H. Leutwyler, *Nucl. Phys.* **B250**, 465 (1985).
- [39] O. Cata and V. Mateu, *J. High Energy Phys.* **09** (2007) 078.
- [40] F.-Z. Chen, M.-D. Zheng, and H.-H. Zhang, *Phys. Rev. D* **106**, 095009 (2022).
- [41] N. Ackermann *et al.*, *Nature (London)* **643**, 1229 (2025).
- [42] D. Z. Freedman, *Phys. Rev. D* **9**, 1389 (1974).
- [43] D. Akimov *et al.* (COHERENT Collaboration), *Science* **357**, 1123 (2017).
- [44] D. Akimov *et al.* (COHERENT Collaboration), *Phys. Rev. Lett.* **126**, 012002 (2021).
- [45] S. Adamski *et al.* (COHERENT Collaboration), *Phys. Rev. Lett.* **134**, 231801 (2025).
- [46] M. Lindner, W. Rodejohann, and X.-J. Xu, *J. High Energy Phys.* **03** (2017) 097.
- [47] W.-F. Chang and J. Liao, *Phys. Rev. D* **102**, 075004 (2020).
- [48] P. Coloma, P. B. Denton, M. C. Gonzalez-Garcia, M. Maltoni, and T. Schwetz, *J. High Energy Phys.* **04** (2017) 116.
- [49] R. H. Helm, *Phys. Rev.* **104**, 1466 (1956).
- [50] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, *Comput. Phys. Commun.* **180**, 747 (2009).
- [51] G. Bélanger, F. Boudjema, A. Goudelis, A. Pukhov, and B. Zaldivar, *Comput. Phys. Commun.* **231**, 173 (2018).
- [52] E. Cortina Gil *et al.* (NA62 Collaboration), *J. High Energy Phys.* **02** (2021) 201.
- [53] Y. M. Andreiev *et al.* (NA64 Collaboration), *Phys. Rev. Lett.* **133**, 121803 (2024).
- [54] T. Feldmann, P. Kroll, and B. Stech, *Phys. Rev. D* **58**, 114006 (1998).
- [55] T. Feldmann, P. Kroll, and B. Stech, *Phys. Lett. B* **449**, 339 (1999).
- [56] T. Feldmann, *Int. J. Mod. Phys. A* **15**, 159 (2000).