

Neutrino self-interaction and weak mixing-angle measurements

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Neutrino self-interaction with a larger “Fermi constant” is often resorted to for understanding various puzzles of our Universe. We point out that a light, neutrinophilic scalar particle ϕ through radiative correction leads to an energy-scale dependence in the neutrino- Z -boson gauge coupling. The driver behind this phenomenon is a large separation between the mass scales of ϕ and additional heavy particles needed for gauge invariance. This is a generic effect insensitive to details of the UV completion. We show that the running can change the $Z\nu\bar{\nu}$ coupling by several percent and affect the measurement of weak mixing angle through neutrino neutral-current processes. We discuss the interplay between the running of the $Z\nu\bar{\nu}$ coupling and $\sin^2\theta_W$ in various experimental observables. It is possible to disentangle the two effects with more than one precise measurement.

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I. INTRODUCTION

The self-interaction of active neutrinos is a blind spot among the tests of the Standard Model. Its strength is allowed to be much bigger than that dictated by the weak interaction if the self-interaction occurs through the exchange of a light, neutrinophilic mediator. At the cosmic frontier, the presence of self-interacting neutrinos appears favorable by recent analysis for improving the fit to data compared to the standard Λ CDM model [1–8]. Neutrino self-interaction is also well motivated for enabling the production of sterile neutrino dark matter via the early Universe oscillation mechanism [9–14] which is otherwise excluded by astrophysical probes [15]. There have been numerous investigations of neutrino self-interaction at various frontiers, including rare decay processes [16–21], accelerator neutrino experiments [22–24], high-energy colliders [25], ultrahigh-energy neutrino telescopes [26–30], and core-collapse supernova [31–33].

Motivated by the above cosmological clues, we point out a novel role of neutrino self-interaction at the precision frontier of testing the electroweak theory. Through radiative corrections, the presence of a light, neutrinophilic scalar can generate a running effect of the $Z\nu\bar{\nu}$ gauge coupling at energy scales below the Z -pole. We show that existing constraints allow this running effect to modify the low-energy neutrino scattering cross sections via neutral current by as large as a few percent. Currently, there are strong

interests and experiments on the horizon that study neutrino-matter interaction as tests of the Standard Model [34–37]. The neutral-current processes involve another percent-level scale-dependent quantity, the weak mixing angle. We discuss how the running of $Z\nu\bar{\nu}$ coupling can impact the extraction of the weak mixing-angle value from experiments and the interplay of different measurements that can be used to distinguish the two running effects.

II. RADIATIVE CORRECTION TO ELECTROWEAK COUPLINGS

We begin with a simplified model commonly used for neutrino self-interaction analysis,

$$\mathcal{L} = \frac{\lambda}{2} \bar{\nu}^c \nu \phi + \text{H.c.}, \quad (1)$$

where ν is the left-handed neutrino chiral field, and ν^c is its antiparticle, the right-handed antineutrino, and both belong to $SU(2)$ doublets in the Standard Model. The scalar field ϕ is a gauge singlet and has no direct coupling to other Standard Model particles. It can be real or complex, and our discussions do not depend on this choice. The neutrino bilinear field carries two units of lepton number and Eq. (1) is the only Lorentz-invariant and renormalizable way for them to couple to a scalar field without introducing right-handed neutrinos.

In this work, we are interested in the mass of ϕ below the weak scale and explore its radiative corrections to the neutrino neutral-current and charged-current couplings. The tree-level electroweak Lagrangian is

$$\mathcal{L} = g_Z^0 \bar{\nu} \gamma^\mu \mathbb{P}_L \nu Z_\mu + g_W^0 \bar{\ell} \gamma^\mu \mathbb{P}_L \ell W_\mu^- + \text{H.c.}, \quad (2)$$

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where $\mathbb{P}_L = (1 - \gamma_5)/2$, $g_Z^0 = g/(2 \cos \theta_W)$, and $g_W^0 = g/\sqrt{2}$. g is the $SU(2)$ gauge coupling, and θ_W is the weak mixing angle.

At the one-loop level, the relevant diagrams are

$$\begin{aligned} & \text{Diagram 1: } \Gamma_{Z \nu \bar{\nu}}(q^2, m_\phi^2) \equiv i g_Z^0 \Gamma_{Z \nu \bar{\nu}}(q^2, m_\phi^2), \\ & \text{Diagram 2: } \Sigma_\nu(p, m_\phi^2) \equiv -i \Sigma_\nu(p, m_\phi^2), \end{aligned} \quad (3)$$

where $\Gamma_{Z \nu \bar{\nu}}$ and Σ_ν are the amputated vertex correction and neutrino self-energy diagrams. q^μ is the momentum running through the Z field. Because ϕ is neutrinophilic, there are no corrections to the $W \ell \nu$ vertex or the charged lepton self-energy. The functions $\Gamma_{Z \nu \bar{\nu}}$ and Σ_ν in Eq. (3) are the only one-loop corrections that arise from model (1). They are equal to

$$\begin{aligned} \Gamma_{Z \nu \bar{\nu}}(q^2, m_\phi^2) &= -\frac{|\lambda|^2}{32\pi^2} \left\{ \frac{2}{\varepsilon} - 1 + 2 \int_0^1 dx \int_0^{1-x} dy \right. \\ &\quad \times \left[\ln \left(\frac{\mu^2}{(1-x-y)m_\phi^2 - xyq^2} \right) \right. \\ &\quad \left. \left. + \frac{xyq^2}{(1-x-y)m_\phi^2 - xyq^2} \right] \right\}, \\ \Sigma_\nu(p, m_\phi^2) &= -\frac{|\lambda|^2}{32\pi^2} \left(\frac{2}{\varepsilon} + \log \frac{\mu^2}{m_\phi^2} + \frac{1}{2} \right) \not{p}, \end{aligned} \quad (4)$$

where $2/\varepsilon = 2/(4-d) - \gamma + \ln(4\pi)$ is used to regularize ultraviolet (UV) divergence in dimensional regularization, and μ is the renormalization scale. In this calculation, we put the external neutrinos on their mass shell.

The corresponding radiative corrections to the electro-weak couplings can be calculated as

$$\begin{aligned} \delta g_Z(q^2, m_\phi^2) &= g_Z^0 \left(\Gamma_{Z \nu \bar{\nu}}(q^2, m_\phi^2) + \frac{\partial \Sigma_\nu(p, m_\phi^2)}{\partial \not{p}} \right), \\ \delta g_W(m_\phi^2) &= g_W^0 \frac{1}{2} \frac{\partial \Sigma_\nu(p, m_\phi^2)}{\partial \not{p}}. \end{aligned} \quad (5)$$

The UV divergences in $\Gamma_{Z \nu \bar{\nu}}$ and $\partial \Sigma_\nu / \partial \not{p}$ add up rather than cancel each other in δg_Z . Unlike regular Yukawa interactions, ϕ in the Lagrangian (1) turns a neutrino into an antineutrino, resulting in an extra minus sign for the $Z - \nu^c$ coupling in the triangle diagram in (3). Similarly, δg_W only receives contribution from $\partial \Sigma_\nu / \partial \not{p}$ and is clearly UV divergent.

III. IMPLICATIONS OF GAUGE INVARIANCE, UV FINITENESS, AND DECOUPLING THEOREM

The UV divergences found above indicate that the model in Eq. (1) cannot be extrapolated to arbitrarily high scales. By assuming ϕ to be a singlet, it fails to respect the Standard Model gauge invariance.

It is not feasible to assign ϕ a charge under $SU(2) \times U(1)$ which would open up the $Z \rightarrow \phi \phi^*$ decay channel for the light ϕ considered here. The latter makes an order-one correction to the invisible Z decay width and is ruled out by the LEP-II experiment [38–40].

A phenomenologically viable UV completion of Eq. (1) was discussed in [17]. In addition to the gauge singlet scalar ϕ , a $SU(2)$ triplet scalar T with hypercharge +1 is introduced. The interacting Lagrangian reads

$$\mathcal{L}_{UV} = y_1 \bar{L}^c T L + y_2 H^T T^\dagger H \phi + \text{H.c.} \quad (6)$$

In the first term, the electric neutral component of the scalar triplet T^0 couples to $\bar{\nu}^c \nu$. After electroweak symmetry breaking, the second term generates a mixing between ϕ and T^0 . Integrating out the triplet yields Eq. (1) and the relation $\lambda \simeq y_1 y_2 v^2 / M_T^2$ where v is the Higgs condensate. In this concrete model, the triplet T field couples to neutrinos and charged leptons at the most fundamental level. It is straightforward to verify that the UV divergences from various vertex and self-energy diagrams indeed cancel completely. The mixing between ϕ and T^0 is an infrared effect and cannot alter physics in the UV or the cancellation.

The notion about UV finiteness in the above example is not a coincidence but the consequence of the Ward identity. It holds on very general grounds. For any gauge-invariant, renormalizable beyond Standard Model (BSM) Lagrangian that can accommodate Eq. (1), the new particles would only contribute to the gauge couplings' beta functions through the vacuum polarization diagrams and the contribution freezes below the heavy particle mass threshold. The rest of radiative corrections to the gauge couplings must sum up to be finite.

This observation can allow one to write down the $Z \nu \bar{\nu}$ and $W \ell \nu$ gauge couplings in a model-independent way. Consider a generic UV completion of Eq. (1) where ϕ is the only new particle below the weak scale. All other BSM particles are heavy, and M is used to denote their mass scale. The above UV finiteness argument ensures that after including both ϕ and heavy loop contributions to g_Z , g_W , all $1/\varepsilon$ terms must cancel away. Along with it, the renormalization scale μ will be replaced by M , which is now a physical quantity.

For low momentum transfers $|q^2| \ll M^2$, the one-loop renormalized gauge couplings take the form

$$g_Z(q^2, m_\phi^2, M^2) = g_Z^0 - \frac{g_Z^0 |\lambda|^2}{32\pi^2} \left\{ \ln \frac{M^2}{m_\phi^2} + 2 \int dx dy \left[\ln \left(\frac{M^2}{(1-x-y)m_\phi^2 - xyq^2} \right) + \frac{xyq^2}{(1-x-y)m_\phi^2 - xyq^2} \right] + c_Z \right\}, \quad (7)$$

and

$$g_W(m_\phi^2, M^2) = g_W^0 - \frac{g_W^0 |\lambda|^2}{64\pi^2} \left(\ln \frac{M^2}{m_\phi^2} + c_W \right). \quad (8)$$

Most crucially for the following discussion, new physics at mass scale M only serves the role for UV finiteness but has no impact on the momentum transfer q^2 dependence of the renormalized gauge couplings. Indeed, the q^2 dependence of Eq. (7) is identical to that from the light ϕ contribution in Eq. (4). This can be understood because heavy new physics cannot propagate at low energies. Their q^2 -dependent contributions manifest as contact operators suppressed by powers of q^2/M^2 , as dictated by the decoupling theorem [41].

The constants c_Z and c_W in Eqs. (7) and (8) are order-one numbers and depend on details of the UV completion. In the $SU(2)$ triplet scalar extension discussed above, the decoupling theorem sets a stronger requirement that not only the UV divergent terms but also finite terms should cancel between the ϕ and T loop contributions if both are heavy. This can be used to fix $c_Z = -3/2$, $c_W = 0$. More details are given in the Appendix.

IV. RUNNING OF $Z\nu\bar{\nu}$ COUPLING BELOW WEAK SCALE

In the Standard Model, the $Z\nu\bar{\nu}$ gauge coupling does not run below the electroweak scale. In contrast, the renormalized g_Z calculated in Eq. (7) has a logarithmic dependence on the momentum transfer q^2 and runs between M and m_ϕ scales. If the scalar ϕ mediating neutrino self-interaction is the only BSM particle lighter than the Z boson, it generates a novel energy scale dependence in g_Z at low energies. This running effect is enabled thanks to the separation of m_ϕ and M mass scales.

The energy-scale dependence of g_Z can be probed with neutrino neutral-current processes that exchanges the Z boson at low energies. The rest of this paper focuses on neutrino-matter scattering experiments where $q^2 < 0$. We define $Q = \sqrt{-q^2}$ to be the experimental energy scale. It is useful to examine the asymptotic behaviors of renormalized g_Z as a function of Q ,

$$\frac{g_Z}{g_Z^0} = 1 - \frac{|\lambda|^2}{32\pi^2} \times \begin{cases} \ln \frac{M^2}{Q^2} + \ln \frac{M^2}{m_\phi^2} + 2 + c_Z, & Q \gg m_\phi \\ 2 \ln \frac{M^2}{m_\phi^2} + \frac{3}{2} + c_Z, & Q \ll m_\phi \end{cases}. \quad (9)$$

For Q varying across large energy scales, we can resum the leading-log contributions and derive the

renormalization group (RG) equation for g_Z ,

$$\frac{dg_Z}{d \ln Q} = \frac{|\lambda|^2}{16\pi^2} g_Z. \quad (10)$$

The anomalous dimension for the g_Z running, $|\lambda|^2/(16\pi^2)$, is a positive number, implying that g_Z decreases for lower energy processes (smaller Q).

To solve for g_Z using Eq. (10), we still need an experimental input at a particular scale. It is a natural choice to use the Z pole. The presence of heavy new physics in the UV completion of Eq. (1), if not far above the electroweak scale, could modify the global fit to electroweak precision observables. Here, we assume such an effect is small and $g_Z(M_Z)$ takes the best-fit value in the Standard Model. This leads to a useful formula that works for $Q \gg m_\phi$,

$$g_Z^2(Q) \simeq g_Z^2(M_Z) \left(\frac{Q}{M_Z} \right)^{|\lambda|^2/8\pi^2}. \quad (11)$$

Figure 1 depicts the energy scale dependence of g_Z^2 in the neutrino self-interaction model considered in this work, for three choices of m_ϕ . The coupling λ is set to unity. The curves are made using the one-loop result, Eq. (7), and they are insensitive to the value of M as long as it is above M_Z . For large $Q \gg m_\phi$, they numerically agree well with the RG-based formula Eq. (11). The running effect ceases after crossing the m_ϕ threshold.

We find that even for m_ϕ above GeV scale, which passes most existing constraints on the neutrinophilic scalar [42], the change in g_Z^2 can be as large as a few percent level.

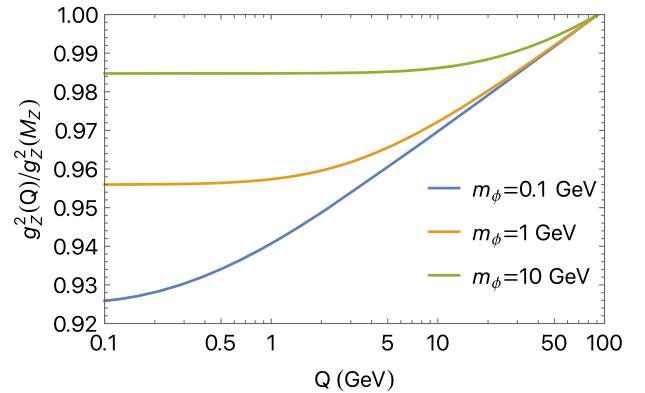


FIG. 1. Running of the $Z\nu\bar{\nu}$ coupling with physical energy scale Q in the presence of a light scalar ϕ that mediates neutrino self-interaction. The coupling λ is set to unity.

The amount of change is comparable to the scale dependence of the weak mixing angle, another running effect to be discussed below.

It is also useful to remark that our discussions so far have ignored the neutrino flavors. In general, the running of g_Z needs not be flavor universal because the neutrino self-interaction introduced in Eq. (1) can be a source of lepton flavor universality violation.

V. IMPACT ON $\sin^2 \theta_W$ MEASUREMENTS WITH NEUTRINOS

The weak mixing angle, or $\sin^2 \theta_W$, is a fundamental parameter of the electroweak theory. At tree level, it is related to the W/Z -boson mass ratio. Including radiative corrections, $\sin^2 \theta_W$ obtains an energy-scale dependence with the leading contribution from loop-generated γ - Z kinetic mixing [43]. Running from the Z pole down to the QCD scale, the Standard Model predicts $\sin^2 \theta_W$ to grow by about 3% [44] which is consistent with most experimental probes.

Existing low-energy limits that resort to electron-hadron and electron-electron interactions will not be affected by the g_Z running found above. However, measurements of $\sin^2 \theta_W$ using neutrino-matter scattering processes will do, because the neutral-current cross sections not only depends on $\sin^2 \theta_W(Q)$ (in the Z -matter coupling) but also is proportional to $g_Z^2(Q)$. If one tries to interpret experimental data by simply assuming the Z -pole value of $g_Z^2(M_Z)$, its running effect would be attributed to that of $\sin^2 \theta_W$, and the result would deviate from the Standard Model prediction. The sign of the shift depends on the target of neutrino scattering. We discuss several cases.

A. Neutrino DIS

The first measurement of $\sin^2 \theta_W$ using neutrino deep-inelastic scattering (DIS) was done by the NuTeV experiment [45]. It shoots a ν_μ or $\bar{\nu}_\mu$ beam onto the target at energy scale $Q \simeq 4.4$ GeV. NuTeV measures the ratios of inclusive neutral-current to charged-current interaction cross sections and found they are lower than the Standard Model predictions (by few percent), leading to a higher $\sin^2 \theta_W$ [46]. While it is commonly appreciated that the NuTeV anomaly could be attributed to hadronic uncertainties [47,48], for the discussions here, it can serve an example to demonstrate the impact of scale dependence in g_Z on the weak mixing-angle measurement. The Q -dependent quantities in the Paschos-Wolfenstein parameter [49] form the combination

$$g_Z^2(Q)[1 - 2\sin^2 \theta_W(Q)]. \quad (12)$$

Using Eq. (11), we find that $g_Z^2(Q)$ runs in the correct direction to explain the deficit in the neutral-current events alone. The orange shaded region in Fig. 2 shows the parameter space that the neutrino self-interaction allows the

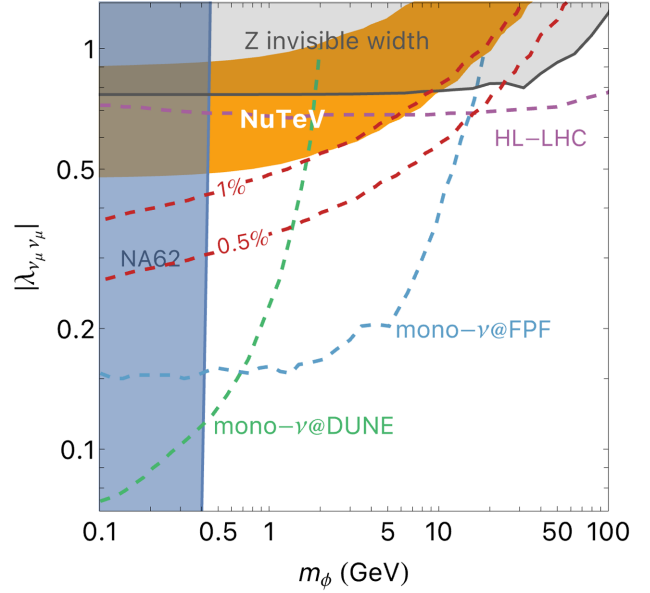


FIG. 2. Parameter space of simplified model with a light, neutrophilic scalar ϕ . The orange shaded region can account for the NuTeV anomaly using the running effect of g_Z while keeping the weak mixing angle consistent with the Standard Model predicted value. The blue-shade region is excluded by the $K \rightarrow \mu\bar{\nu} + \text{invisible}$ measurement by NA62. The gray-shaded region is excluded by invisible Z width measurement by LEP-II. The green-, blue-, and magenta-dashed curves show recent proposals to probe the model by exploring missing energy events at the DUNE near detector, the Forward Physics Facility, and ATLAS/CMS detectors at high-luminosity LHC, respectively. The red-dashed curves correspond to a departure of $g_Z^2(Q)$ from the Z -pole value by 1%, 0.5%, with $Q \ll m_\phi$, and the regions above can be covered if the neutral-current cross section and running of g_Z^2 are measured with percent-level precision at future neutrino experiments.

model in Eq. (1) to be consistent with NuTeV within 2σ level (we use data from Fig. 1 of [44]), while still keeping $\sin^2 \theta_W(Q)$ as predicted in the Standard Model. We also show a recent constraint from exotic kaon decay measurement by the NA62 experiment [50], which sets a lower limit on the viable ϕ mass, and from the invisible Z -decay width measurement by LEP-II [51,52]. The green-, blue-, and magenta-dashed curves show recent proposals to cover the parameter space by exploring missing energy events at the DUNE near detector [22], the Forward Physics Facility [23], and ATLAS/CMS detectors at high-luminosity LHC [25], respectively. The novel Q dependence of neutrino DIS [Eq. (12)] can be further examined with the scattering of LHC neutrinos at the Forward Physics Facility [53].

B. CE ν NS

The coherent elastic neutrino-nucleus scattering (CE ν NS) is another channel to probe the running of g_Z . The CE ν NS process has already been detected by COHERENT [54], XENONnT [55,56], and PandaX-4T

[57]. More experiments are on the horizon to explore the CE ν NS process with higher precision [58,59]. The CE ν NS cross section depends on g_Z and $\sin^2 \theta_W$ as [60,61]

$$g_Z^2(Q)[N - Z(1 - 4\sin^2 \theta_W(Q))]^2, \quad (13)$$

where the square bracket is nuclear weak charge. Assuming an isoscalar target with equal number of protons and neutrons ($N = Z$), the cross section is proportional to

$$g_Z^2(Q) \sin^4 \theta_W(Q). \quad (14)$$

The decrease of g_Z^2 at low energies could be interpreted as a smaller weak mixing angle than the Standard Model value at the corresponding energy scale. Interestingly, this is an opposite effect to that for neutrino DIS.

C. $\nu_\mu(\bar{\nu}_\mu)$ -electron scattering

Recent works make projections for measuring the weak mixing angle at accelerator neutrino experiments like DUNE and SBND with an electron target [62,63]. These experiments can be tuned to run with either a neutrino or antineutrino beam. The $\nu_\mu(\bar{\nu}_\mu)$ -electron scattering cross sections have different parametrical dependencies. In the $E_\nu \gg m_e$ limit,

$$\begin{aligned} \sigma_{\nu_\mu e \rightarrow \nu_\mu e} &\propto g_Z^2(Q)[3 - 12\sin^2 \theta_W(Q) + 16\sin^4 \theta_W(Q)], \\ \sigma_{\bar{\nu}_\mu e \rightarrow \bar{\nu}_\mu e} &\propto g_Z^2(Q)[1 - 4\sin^2 \theta_W(Q) + 16\sin^4 \theta_W(Q)]. \end{aligned} \quad (15)$$

For $\sin^2 \theta_W$ around 0.23, the first square bracket is a decreasing function of $\sin^2 \theta_W$, whereas the second one is an increasing function. If neutrino self-interaction is at work and reduces $g_Z^2(Q)$ at low energies, it will lead to a suppression in both $\sigma_{\nu_\mu e \rightarrow \nu_\mu e}$ and $\sigma_{\bar{\nu}_\mu e \rightarrow \bar{\nu}_\mu e}$. If one tried to attribute this to the running of the weak mixing angle, they would reach opposite conclusions by using neutrino and antineutrino beams—the former channel calls for a higher $\sin^2 \theta_W(Q)$ but the latter lower.

All the neutrino scattering cross sections discussed above are proportional to g_Z^2 , but they depend on $\sin^2 \theta_W$ in different ways. With more than one precise measurement, one will be able to tell their running effects apart. Furthermore, if future neutrino neutral-current measurements are sensitive to the percent-level g_Z^2 running effect from the Z pole down to low energies, they can cover the regions above the red-dashed curves in Fig. 2.

VI. CONCLUSION

Neutrino self-interaction via a light mediator particle is of great interest to address puzzles from the sky. We work on a simple extension of the Standard Model with a neutrinophilic scalar ϕ and calculate its radiative correction to the electroweak couplings. We derive a novel

energy-scale dependence in the value of $Z\nu\bar{\nu}$ coupling below the Z -boson mass. This effect is granted by a large separation of m_ϕ and mass scale of other heavy particles in gauge-invariant UV completions. Thanks to the elusive nature of neutrinos, their couplings to the neutrinophilic mediator are still allowed to be sizable, and the corresponding running effect of the $Z\nu\bar{\nu}$ coupling is allowed to be as large as several percent. It affects the neutrino scattering cross sections via neutral-current interactions and in turn measurements of the weak mixing angle, another scale-dependent quantity. We discuss the interplay of the two running effects and point out that they can be differentiated by combining measurements. There is rich opportunity to disclose novel neutrino self-interaction at upcoming neutrino experiments.

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APPENDIX

In this appendix, we derive the constants c_Z, c_W in the renormalized gauge couplings [Eqs. (7) and (8)], for the model with a $SU(2)$ triplet scalar [Eq. (6)] as the UV completion of Eq. (1). The discussion is based on the decoupling theorem.

First, we set the parameter $y_2 = 0$. In this case, the mixing between ϕ and T^0 vanishes, and ϕ completely decouples from other particles. All the BSM radiative corrections are from the interaction between T and the lepton doublet L . If we give the whole $SU(2)$ triplet a gauge-invariant and large mass M , all its effects on the Standard Model as the low-energy effective theory must be suppressed by inverse powers of M .

Next, we turn on a nonzero y_2 and in turn the ϕ - T^0 mixing. This will reshuffle radiative corrections (and the UV divergences, which will still cancel) to various loop diagrams involving ϕ and T . The mixing will enter through various vertex Feynman rules. To proceed along the above argument, we are after the case where the ϕ - T^0 mixing does not vanish in the decoupling limit. It is possible to realize this situation by examining the ϕ - T^0 mixing mass matrix,

$$\begin{pmatrix} \mu_T^2 & y_2 v^2 \\ y_2 v^2 & \mu_\phi^2 \end{pmatrix}, \quad (A1)$$

where μ_T^2 and μ_ϕ^2 are from the quadratic T and ϕ potential terms before the mixing. The useful decoupling limit for our discussion here will be sending both μ_T and μ_ϕ to infinity while keeping them equal to each other. In this case,

ϕ and T^0 are nearly degenerate, and the mass matrix is always diagonalized with a $(\pi/4)$ mixing angle. As a result, the induced ϕ -neutrino coupling does not vanish, neither do the diagrams in Eq. (3).

By sending both ϕ and T mass to infinity, the decoupling of their joint effects should continue to work. It means if one includes all the radiative correction contributions from ϕ and T , not only their UV divergence but also the finite terms must cancel away, up to q^2/M^2 terms. This argument provides a simple way to evaluate the (many) T -loop diagrams without needing to directly compute them, because we have already calculated the ϕ ones for general ϕ mass. The heavy T contributions to g_Z and g_W are given by

$$\begin{aligned}\delta g_Z^{(T)} &= - \lim_{M^2 \gg |q^2|} \delta g_Z(q^2, M^2) \\ &= \frac{|\lambda|^2}{16\pi^2} \left(\frac{2}{\epsilon} - \ln \frac{\mu^2}{M^2} + \frac{1}{2} \right), \\ \delta g_W^{(T)} &= - \lim_{M^2 \gg |q^2|} \delta g_W(M^2) \\ &= \frac{|\lambda|^2}{64\pi^2} \left(\frac{2}{\epsilon} - \ln \frac{\mu^2}{M^2} + \frac{1}{2} \right),\end{aligned}\quad (\text{A2})$$

where the functions $\delta g_{Z,W}$ are given by Eq. (5) and M denotes the physical mass of T .

By dialing the new physics parameters, it is possible to make ϕ light while keeping T still heavy. The ϕ, T^0 couplings will change correspondingly, but the number of Feynman diagrams relevant for radiative corrections and the structure of the loop functions (q^2 and mass dependence) still hold as before. As long as the mass of T remains well above the momentum transfer, the above forms of $\delta g_{Z,W}^{(T)}$ will hold intact. Adding them to the ϕ contributions yields Eqs. (7) and (8) with $c_Z = -3/2$, $c_W = 0$.

A consistency check is the $Q \ll m_\phi$ limit of Eq. (9). With $c_Z = -3/2$, we find the radiative correction to g_Z is proportional to $\ln(M^2/m_\phi^2)$. With $c_W = 0$, the radiative correction of g_W in Eq. (8) is proportional to the same logarithmic factor. In the $(\pi/4)$ mixing scenario discussed above (with $\mu_T = \mu_\phi$), Eq. (A1) gives $M^2 = \mu_T^2 + y_2 v^2$ and $m_\phi^2 = \mu_T^2 - y_2 v^2$. The logarithmic factor vanishes as $\sim v^2/\mu_T^2$ for very large μ_T , fulfilling the desired decoupling behavior.

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