

**Minimal dark  $SU(2)$  origin of a massless Dirac neutrino**P. S. Bhupal Dev<sup>1,2</sup>, Julia Gehrlein<sup>3</sup>, Amartya Sengupta<sup>4,5</sup> and Amarjit Soni<sup>6</sup><sup>1</sup>*Department of Physics and McDonnell Center for the Space Sciences,  
Washington University, Saint Louis, Missouri 63130, USA*<sup>2</sup>*PRISMA<sup>++</sup> Cluster of Excellence and Mainz Institute for Theoretical Physics,  
Johannes Gutenberg-Universität Mainz, 55099 Mainz, Germany*<sup>3</sup>*Department of Physics, Colorado State University, Fort Collins, Colorado 80523, USA*<sup>4</sup>*Department of Physics, The State University of New York, Buffalo, New York 14260, USA*<sup>5</sup>*International Center for Quantum-field Measurement Systems for Studies of the Universe and Particles  
(QUP), KEK, 1-1 Oho, Tsukuba, Ibaraki 305-0801, Japan*<sup>6</sup>*Physics Department, High Energy Theory Group, Brookhaven National Laboratory,  
Upton, New York 11973, USA*

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We propose a gauge-symmetry origin of a rank-two Dirac neutrino mass matrix that enforces one exactly massless neutrino, while being consistent with the oscillation data, as well as cosmological constraints. The mechanism relies on a minimal dark  $SU(2)_D$  gauge symmetry under which one right-handed neutrinolike Weyl fermion is charged, thereby forbidding its Standard Model Yukawa coupling. Quantum consistency then fixes the minimal dark-sector completion: Cancellation of the Witten anomaly requires a second fermionic  $SU(2)_D$  doublet, while a discrete  $Z_4$  symmetry that forbids Majorana masses allows the two dark doublets to form a vectorlike pair. This anomaly-free completion gives rise to a secluded, confining dark sector that can contain a potential dark matter candidate, linking the protected neutrino texture to dark infrared dynamics.

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**Introduction.** The observation of neutrino oscillations [1,2] has increased the number of free parameters in the Standard Model (SM) by at least seven: three mixing angles, three neutrino masses, and at least one  $CP$  phase. While neutrino oscillation experiments are making steady progress in measuring the oscillation parameters more precisely [3], with further improvements anticipated from ongoing and next-generation experiments like JUNO [4], IceCube/DeepCore [5], Hyper-Kamiokande [6], DUNE [7], and their combinations [8–12], none of the oscillation experiments are sensitive to the absolute neutrino mass scale. The observed solar and atmospheric mass-squared splittings require two nonzero neutrino masses, while the mass of the lightest neutrino cannot be constrained by oscillation experiments and remains an open question. Instead, we have to rely on cosmological observations of large-scale structure, which constrain the sum of neutrino masses [13], and precise laboratory measurements of the end point of the  $\beta$ -decay spectrum [14] to provide complementary information on the neutrino mass scale. The direct laboratory searches, while impressive, are currently far from the sensitivity required to probe even the inverted ordering

(IO), let alone normal ordering (NO) [15]. Additional information on the neutrino masses can be obtained if a Galactic supernova is detected [6,16–26] or a measurement of the cosmic neutrino background is achieved [27,28].

Currently, the strongest constraints on the neutrino mass scale come from the recent DESI baryon acoustic oscillation (BAO) analysis [29]. When interpreted within the  $\Lambda$  cold dark matter ( $\Lambda$ CDM) cosmology [30,31], it places an upper bound on the sum of neutrino masses  $\sum_i m_i < 64$  meV [95% confidence level (CL)]. Slightly weaker constraints are obtained from Planck [30], ACT [32] and SPT [33] datasets, while a combination of ACT + SPT + Planck gives a slightly stronger bound  $\sum_i m_i < 62$  meV (95% CL) [34]. These upper bounds are close to the minimum value of  $\sum m_i$  allowed by oscillation data for NO and already in tension with the corresponding value for IO. Any emerging tension with NO can be alleviated if the lightest neutrino is (nearly) massless. Although this conclusion might change depending on the assumed cosmological model [35–53] and the presence of additional neutrino interactions [54–59], the increasing precision of cosmological constraints provides a timely motivation to consider theories in which the lightest neutrino remains exactly massless.<sup>1</sup>

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<sup>1</sup>Additional theoretical motivation for a massless neutrino comes from models with a  $CPT$ -symmetric Universe [60,61].

Existing models with one massless neutrino rely on a minimal field content of just two right-handed neutrinos instead of the three needed to generate masses for all SM neutrinos. Popular ones include the minimal scotogenic model [62–65], minimal type-I seesaw [66–69] constructions with two right-handed neutrinos [70–78], and variants of the type-I seesaw such as inverse-seesaw- [79] based minimal models [80–84]. These models are phenomenologically viable, although if one tries to gauge the usual  $B - L$  symmetry with only two ordinary right-handed neutrinos, additional chiral states or nonstandard charge assignments are required for anomaly cancellation. This prevents these models from featuring an additional gauged  $B - L$  symmetry, which is desirable to protect the neutrino mass, and also for generating a massive gauge boson after lepton-number breaking associated with Majorana mass terms or embedding it within grand unified theories where  $B - L$  is part of a bigger gauge group. An alternative approach is to postulate a texture of the neutrino mass matrix which leads to only two nonzero eigenvalues [85–88]. However, in this case the zero neutrino mass is, in general, not protected from higher-order corrections, rendering it nonzero at a certain energy scale [89], although in some models the texture zeros are stable [90,91].

Here, we instead show that a massless neutrino can follow from a gauge selection rule. Because of its symmetry origin, the zero neutrino mass is protected from corrections and one neutrino remains massless at all energy scales and all orders in perturbation theory. As a minimal example, we introduce a  $SU(2)_D$  gauge symmetry under which one right-handed neutrino-like SM singlet Weyl fermion  $N_1$  transforms as a doublet. Furthermore, quantum consistency constrains the additional dark fermion content, which can accommodate a dark matter (DM) candidate, as shown below.<sup>2</sup>

*The minimal model.* The minimal choice of the dark gauge group is  $SU(2)_D$  which features three massless dark gauge fields that couple to the fermions charged under this group with coupling strength  $g_D$ . To obtain an anomaly-safe dark gauge theory, the global Witten anomaly needs to be prevented, which requires the introduction of an even number of  $SU(2)_D$  doublets [94]. Therefore, we include, in addition to  $N_1$ , another  $SU(2)_D$  doublet Weyl-fermion field  $\xi$ . The SM fields and two right-handed neutrinos  $N_2$ ,  $N_3$  are singlets under the dark gauge group. All new fermions are singlets under the SM gauge groups. Additionally, we impose a discrete  $Z_4$  symmetry, acting on the SM fermions and on the new singlet and dark-sector

fermions, to keep the neutrino sector Dirac. The masslessness of one neutrino is enforced by the dark gauge charge, which forbids one Dirac Yukawa column. However, the dark gauge symmetry alone does not exclude Majorana-type bilinears for the SM singlet fermions; if present, these terms would turn the model into a Majorana or seesawlike construction rather than a rank-two Dirac neutrino model. The chosen  $Z_4$  assignment allows the charged-lepton and Dirac neutrino Yukawa terms and forbids Majorana masses, but permits the mixed dark bilinear involving  $N_1$  and  $\xi$ . A  $Z_2$  symmetry acting on the same fermions would not be sufficient, since the Majorana bilinears would be even under it. The field content and their charges are shown in Table I.

With these charge assignments, the allowed renormalizable Yukawa interactions with the SM lepton and Higgs sector are

$$\mathcal{L}_Y = -(Y_e)_{ij} \bar{L}_i H e_{Rj} - (Y_\nu)_{i\alpha} \bar{L}_i \tilde{H} N_\alpha + \text{H.c.}, \quad (1)$$

where  $\tilde{H} = i\sigma_2 H^*$ ;  $i, j = 1, 2, 3$ , and  $\alpha = 2, 3$ . The Majorana-type operators

$$N_\alpha N_\beta, \quad \epsilon_{ab} N_1^a N_1^b, \quad \epsilon_{ab} \xi^a \xi^b \quad (2)$$

carry charge two modulo four and are therefore absent in the  $Z_4$ -symmetric theory. Here  $\epsilon_{ab}$ , with  $\epsilon_{12} = +1$ , denotes the antisymmetric  $SU(2)_D$ -invariant tensor used to contract two dark-doublet indices. The mixed dark bilinear  $\epsilon_{ab} N_1^a \xi^b$  has charge  $1 + 3 = 4 = 0 \bmod 4$ , so the vectorlike mass

$$\mathcal{L}_{\text{mass}} = -m_D \epsilon_{ab} N_1^a \xi^b + \text{H.c.} \quad (3)$$

is allowed and it controls the dark infrared dynamics that feature a confinement scale  $\Lambda_D$ .

For  $m_D \gg \Lambda_D$ , the dark doublets decouple and the low-energy theory approaches pure  $SU(2)_D$  Yang-Mills theory; for  $m_D \lesssim \Lambda_D$ , the dark fermions remain part of the confined spectrum. The corresponding dark-sector phenomenology depends on this hierarchy, as discussed later.

The protected zero column in the neutrino mass matrix is not filled by higher-dimensional operators as long as the  $SU(2)_D$  and  $Z_4$  symmetries remain exact. Any operator that

TABLE I. Charges of the SM leptons, the Higgs field, and the dark-sector fermions under  $SU(2)_D$  and  $Z_4$ .

| Field      | $SU(3)_c \times SU(2)_L \times U(1)_Y$ | $SU(2)_D$    | $Z_4$ |
|------------|----------------------------------------|--------------|-------|
| $L_i$      | $(\mathbf{1}, \mathbf{2}, -1/2)$       | $\mathbf{1}$ | 1     |
| $e_{Ri}$   | $(\mathbf{1}, \mathbf{1}, -1)$         | $\mathbf{1}$ | 1     |
| $H$        | $(\mathbf{1}, \mathbf{2}, 1/2)$        | $\mathbf{1}$ | 0     |
| $N_1$      | $(\mathbf{1}, \mathbf{1}, 0)$          | $\mathbf{2}$ | 1     |
| $\xi$      | $(\mathbf{1}, \mathbf{1}, 0)$          | $\mathbf{2}$ | 3     |
| $N_2, N_3$ | $(\mathbf{1}, \mathbf{1}, 0)$          | $\mathbf{1}$ | 1     |

<sup>2</sup> $SU(2)_D$  in the context of Dirac neutrino mass and DM has also been considered in Refs. [92,93], where the  $SU(2)_D$  gauge symmetry is spontaneously broken to retain a global  $SU(2)$  symmetry. Two Dirac neutrino masses are generated at one loop from the  $SU(2)_D$  dark fermions and scalars. In our mechanism, the  $SU(2)_D$  gauge symmetry remains unbroken and Dirac neutrino masses arise at tree level.

would generate the forbidden Yukawa coupling  $\bar{L}_i \tilde{H} N_1$  carries an uncontracted  $SU(2)_D$  fundamental index unless it is multiplied by a dark-sector operator transforming as a fundamental of  $SU(2)_D$ . Since the  $SU(2)_D$  symmetry is unbroken, and confinement does not break gauge invariance, only  $SU(2)_D$ -singlet operators can acquire vacuum expectation values. Thus, dark confinement can generate singlet bound states or condensates, but it cannot induce the forbidden Yukawa column. The  $Z_4$  symmetry separately forbids Majorana-type operators, including higher-dimensional lepton-number-violating terms, so the active Dirac neutrino mass matrix remains rank two.

*Minimality and alternative choices of the dark symmetry.* The zero-column mechanism itself does not require a non-Abelian group; an Abelian dark charge could also forbid  $\bar{L}_i \tilde{H} N_1$ . We choose a non-Abelian realization because it also produces a secluded confining sector with an infrared scale generated by strong dynamics. With this added requirement,  $SU(2)_D$  is the smallest possible choice. Its minimal consistency issue is the Witten anomaly, already removed by the second doublet  $\xi$  introduced above.

A hidden  $SU(3)_D$  theory is a natural alternative. It is similar to QCD and can support dark baryons, but a single Weyl fundamental of  $SU(3)_D$  is perturbatively anomalous. The smallest anomaly-free version introduces an antifundamental partner,  $\psi_D \sim \mathbf{3}$ ,  $\eta_D \sim \bar{\mathbf{3}}$ , making the dark sector vectorlike. Such models may be preferable for dark baryon stability or first-order confinement, but they are less minimal to predict a massless neutrino. Other non-Abelian groups like  $SO(N \geq 4)$  and  $Sp(N \geq 2)$  can also realize the zero-column mechanism, but their spectra and accidental symmetries differ from the minimal  $SU(2)_D$  case.

*Neutrino mass matrix.* After electroweak symmetry breaking, the Lagrangian (1) generates a Dirac neutrino mass matrix

$$M_\nu = \frac{v}{\sqrt{2}} Y_\nu, \quad Y_\nu = \begin{pmatrix} 0 & y_{12} & y_{13} \\ 0 & y_{22} & y_{23} \\ 0 & y_{32} & y_{33} \end{pmatrix}, \quad (4)$$

where  $v$  is the electroweak vacuum expectation value. The first column of  $Y_\nu$  is identically zero because  $N_1$  has a dark gauge charge which does not allow its coupling to the SM neutrinos. We nevertheless display the texture in a  $3 \times 3$  square form by reserving the first column for the dark-charged state  $N_1$ . Hence,  $\det M_\nu = 0$ , and one physical Dirac neutrino mass vanishes. Working in the charged-lepton mass basis, the Dirac mass matrix is diagonalized by

$$U_{\text{PMNS}}^\dagger M_\nu U_R = \text{diag}(m_1, m_2, m_3), \quad (5)$$

where  $U_R$  is an arbitrary unitary matrix diagonalizing the  $N_i$  sector and  $U_{\text{PMNS}}$  is the measurable Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix [95,96].

Furthermore, since  $N_2$  and  $N_3$  carry identical charges, we can perform a right-handed  $U(2)$  rotation among them,<sup>3</sup>

$$\begin{pmatrix} N_2 \\ N_3 \end{pmatrix} \rightarrow R \begin{pmatrix} N_2 \\ N_3 \end{pmatrix}, \quad R \in U(2). \quad (6)$$

Let  $B$  be the  $3 \times 2$  matrix formed by the nonzero columns of  $M_\nu$ . If the second row of  $B$  is  $(w_1, w_2)$ , choose

$$R^\dagger = \frac{1}{\sqrt{|w_1|^2 + |w_2|^2}} \begin{pmatrix} w_2 & w_1^* \\ -w_1 & w_2^* \end{pmatrix}. \quad (7)$$

Then the second-row entry of the rotated two-column block vanishes, and the full matrix can be written as

$$M_\nu = \begin{pmatrix} 0 & m_{12} & m_{13} \\ 0 & 0 & m_{23} \\ 0 & m_{32} & m_{33} \end{pmatrix}. \quad (8)$$

This is the  $D_{76}$  class of four-zero texture Dirac mass matrix [88].<sup>4</sup> In the present model, this is a basis representation, not an additional flavor prediction. Since the two nonzero columns correspond to fields with identical charges, the additional zero can equivalently be placed in another entry of the second or third column by an appropriate right-handed  $U(2)$  basis choice. These choices correspond to other  $D$ -type textures with one zero column in the texture-zero classification, but only  $D_{76}$  in Eq. (8) and the electron-row representative  $D_{16}$  (with  $m_{12} = 0$ ) are compatible with the current oscillation data for both NO and IO [88]. We focus on the flavor completion of  $D_{76}$  because the physical zero  $(M_\nu)_{\mu 2} = 0$  involves  $\mu$ -row PMNS elements, which depend on both  $\theta_{23}$  and  $\delta_{CP}$ :  $U_{\mu 2}$ ,  $U_{\mu 3}$  for NO and  $U_{\mu 1}$ ,  $U_{\mu 2}$  for IO, and can therefore correlate  $\theta_{23}$  with  $\delta_{CP}$  once the right-handed basis is fixed by additional flavor structure (see Supplemental Material [99]). We reiterate that the only basis-independent prediction of the minimal  $SU(2)_D$  construction is the protected zero column and, hence, one massless Dirac neutrino; the  $D_{76}$  texture and the resulting  $\theta_{23}$ – $\delta_{CP}$  correlations are predictions only of a flavor completion that fixes the right-handed basis.

*Phenomenology.* Neutrino sector: As stated before, we predict a zero lightest neutrino mass in this model. The other neutrino masses are fixed by the measured mass splittings  $\Delta m_{21}^2$  and  $\Delta m_{31}^2$ ,

<sup>3</sup>This freedom of choosing basis in the  $(N_2, N_3)$  space follows solely from the identical quantum numbers of  $N_2$  and  $N_3$  and is independent of the particular choice of dark symmetry group.

<sup>4</sup>Four-zero texture neutrino masses have been discussed in Ref. [97]. An  $SO(10)$  grand unified theory embedding of related texture-zero structures has been studied in Ref. [98].

$$\begin{aligned}
\text{NO: } m_1 &= 0, \quad m_2 = \sqrt{\Delta m_{21}^2}, \quad m_3 = \sqrt{\Delta m_{31}^2}, \\
\text{IO: } m_3 &= 0, \quad m_1 = \sqrt{|\Delta m_{31}^2|}, \quad m_2 = \sqrt{|\Delta m_{31}^2| + \Delta m_{21}^2}.
\end{aligned}
\tag{9}$$

The observable related to the end point of the  $\beta$ -decay spectrum is [15]

$$m_\beta^2 = \sum_i |U_{ei}|^2 m_i^2, \tag{10}$$

with  $U_{ei}$  the PMNS matrix elements in the electron row. Using the recent JUNO results for the solar sector [104], together with the NUFIT6.0 best-fit values for the remaining oscillation parameters [105], we find for Eq. (10) and the sum of the neutrino masses

$$\begin{aligned}
\text{NO: } m_\beta &= 8.85_{-0.57}^{+0.59} \text{ meV}, \quad \sum_i m_i = 58.79_{-1.04}^{+1.07} \text{ meV}, \\
\text{IO: } m_\beta &= 48.76_{-0.88}^{+1.08} \text{ meV}, \quad \sum_i m_i = 98.92_{-1.65}^{+2.06} \text{ meV},
\end{aligned}
\tag{11}$$

where the error bars account for the  $3\sigma$  oscillation-parameter uncertainties (see Fig. 1). The predicted values of  $m_\beta$  are far below the current KATRIN bound,  $m_\beta < 450$  meV [15]. For the sum of neutrino masses, the NO prediction lies below both the standard Planck + BAO  $\Lambda$ CDM bound  $\sum_i m_i < 120$  meV [30] and the stronger DESI + cosmic microwave background (CMB)  $\Lambda$ CDM bound  $\sum_i m_i < 64$  meV [29]. The IO prediction remains allowed by Planck + BAO but disfavored by DESI + CMB.

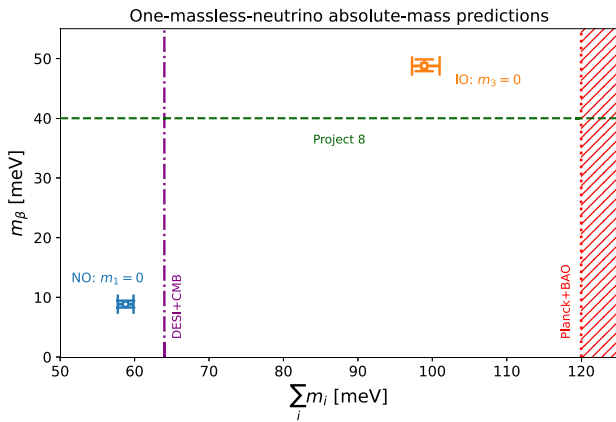


FIG. 1. Absolute-mass predictions for the zero-column representative texture with one massless Dirac neutrino. Error bars show the propagated uncertainties obtained by scanning representative  $3\sigma$  oscillation-parameter ranges in both  $\sum_i m_i$  and  $m_\beta$ . The vertical lines indicate the DESI + CMB [29] and Planck + BAO [30]  $\Lambda$ CDM reference bounds on  $\sum_i m_i$ , and the green dashed line shows the approximate Project 8 target sensitivity near  $m_\beta \simeq 40$  meV [106]. The current KATRIN bound on  $m_\beta$  lies above the plotted range [15].

Future improvements in direct  $\beta$ -decay searches from experiments like KATRIN [107], ECHO [108], HOLMES [109] and Project 8 [106] could eventually reach sensitivities down to  $m_\beta \simeq 40$  meV and probe the IO case, while the NO prediction is more challenging to reach. Future CMB and large-scale-structure surveys can further test the mass-sum prediction, with CMB-S4 + DESI-like combinations expected to reach  $\mathcal{O}(10 \text{ meV})$  precision in representative forecasts [110–113]. This would directly probe the NO prediction near 59 meV. Figure 1 summarizes the relevant absolute-mass probes.

This model features Dirac neutrinos, which additionally predicts that neutrinoless double  $\beta$  decay should not be observed. Also, the relic neutrino capture rate is smaller by a factor of 2 compared to the Majorana case, giving about 4.06 events/yr for a 100 g PTOLEMY-like detector [114]. For a detector with an arbitrarily good energy resolution, each mass eigenstate  $\nu_i$  would make a distinguishable contribution to the C $\nu$ B capture and to the  $\beta$ -decay spectrum. The  $\beta$ -decay spectrum would be the sum of three spectra, and its end point would be determined by the lightest neutrino mass, which offers a way to test whether the lightest mass is zero or not. However, in practice, the primary target energy resolution of PTOLEMY is  $\sim 50$  meV [27], so, experimentally, it will be extremely difficult, if not impossible, to determine whether the lightest neutrino is exactly massless or has a nonzero mass smaller than the detector resolution.

**Dark sector:** The neutrino results derived in the previous subsection do not require knowledge of the detailed spectrum of the dark gauge theory; instead, it follows from the nontrivial dark charge of only  $N_1$ . The dark-sector phenomenology, however, depends on the choice of the gauge group and the particle spectrum.

Once  $SU(2)_D$  is gauged, the theory becomes strong in the infrared and generates its own mass scale by dimensional transmutation [115,116],

$$\Lambda_D = \mu \exp \left[ -\frac{8\pi^2}{b_0 g_D^2(\mu)} \right]. \tag{12}$$

Here  $b_0$  is the one-loop  $\beta$ -function coefficient. With  $n_W$  Weyl-fermion doublets and  $n_S$  complex scalar doublets of  $SU(2)_D$ ,

$$b_0 = \frac{22}{3} - \frac{n_W}{3} - \frac{n_S}{6}. \tag{13}$$

In the minimal model  $n_W = 2$  and  $n_S = 0$ , so  $b_0 = 20/3 > 0$ ; the dark gauge theory is therefore asymptotically free and is expected to confine in the infrared.<sup>5</sup>

<sup>5</sup>Losing asymptotic freedom at one loop would require  $2n_W + n_S \geq 44$ , which gives  $n_W \geq 22$  Weyl doublets in the absence of scalar doublets.

As in QCD, a gauge theory with no elementary mass scale in the gauge sector can produce massive bound states after confinement. If the lightest confined states are stable or sufficiently long-lived, the particles in such a secluded confining sector provide natural DM candidates [117–130].

The confined spectrum depends on the hierarchy between  $m_D$ , the mass of the dark fermions, and the confinement scale  $\Lambda_D$ . If  $m_D \gg \Lambda_D$ , the dark doublets are heavy compared to the confinement scale, and the low-energy theory is close to pure Yang-Mills theory. The lightest states are then dark glueballs. The lightest scalar and pseudoscalar states may be written schematically in terms of the dark field strength tensor  $G_{\mu\nu}^A$  and its dual  $\tilde{G}^{A\mu\nu}$  as

$$S_D \sim G_{\mu\nu}^A G^{A\mu\nu}, \quad P_D \sim G_{\mu\nu}^A \tilde{G}^{A\mu\nu}, \quad (14)$$

where  $A = 1, 2, 3$  is the adjoint  $SU(2)_D$  gauge index. Their masses are of order

$$m_{S_D}, m_{P_D} \sim \kappa \Lambda_D, \quad (15)$$

where  $\kappa$  denotes nonperturbative coefficients derived from lattice studies of confining Yang-Mills spectra [131,132]. If instead  $m_D \lesssim \Lambda_D$ , the dark doublets participate in the confined spectrum. The theory can then contain fermionic and bosonic dark hadrons built from  $N_1$  and  $\xi$ ; for example, states like

$$\epsilon_{ab} N_1^a \xi^b, \quad N_1^\dagger \bar{\sigma}^\mu \xi \quad (16)$$

arise, where  $\bar{\sigma}^\mu = (1, -\vec{\sigma})$  in two-component Weyl notation and  $\vec{\sigma}$  are the Pauli matrices. The detailed spectrum is nonperturbative, but the existence of a secluded confining sector follows directly from the anomaly-safe completion.

Whether the lightest confined state in our model is a viable DM candidate depends on its stability, portal interactions, relic abundance, and dark-sector thermal history. In the discussion below, the DM interpretation assumes that the lightest confined state is protected by an accidental dark-sector symmetry or that its decay rate is sufficiently suppressed to make it cosmologically long-lived. In the pure Yang-Mills limit,  $m_D \gg \Lambda_D$ , the lightest dark glueball is stable if no portal to the SM is present and becomes long-lived when its decay proceeds only through higher-dimensional operators coupling the Higgs field to dark glue. If  $m_D \lesssim \Lambda_D$ , the lightest dark hadron may instead contain the dark fermions  $N_1, \xi$ . Its stability then depends on the accidental global symmetries of the renormalizable dark-sector Lagrangian. In particular, the gauge interactions and the vectorlike mass term can preserve a dark fermion number, analogous to baryon number in QCD, even if this symmetry is not imposed explicitly. Higher-dimensional portal operators or additional ultraviolet (UV) interactions may violate this accidental symmetry and determine the lifetime of the lightest dark hadron.

The leading interactions of the dark sector with the SM arise through higher-dimensional portals. Since the dark gauge group is non-Abelian, there is no dimension-four kinetic mixing with hypercharge. The dark gauge theory can, however, contain its own topological term,

$$\mathcal{L}_{\theta_D} = \frac{\theta_D g_D^2}{32\pi^2} G_{\mu\nu}^A \tilde{G}^{A\mu\nu}, \quad \tilde{G}^{A\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} G_{\rho\sigma}^A. \quad (17)$$

This term is internal to the dark sector and is not a portal to the SM. Its physical meaning depends on the dark-fermion mass structure. In the chiral limit, when an anomalous axial rotation of the dark fermions is available,  $\theta_D$  can be rotated away. For nonzero  $m_D$ , the physical  $CP$ -odd parameter is the dark analog of  $\bar{\theta}$ , schematically  $\bar{\theta}_D = \theta_D + \arg m_D$ , up to convention-dependent signs.<sup>6</sup> It may affect dark  $CP$  violation, the confined spectrum, or the finite-temperature transition, but it does not by itself lead to visible-sector interactions. The simplest Higgs portals involving dark gauge fields appear at dimension six,

$$\frac{c_G}{\Lambda^2} (H^\dagger H) G_{\mu\nu}^A G^{A\mu\nu}, \quad \frac{\tilde{c}_G}{\Lambda^2} (H^\dagger H) G_{\mu\nu}^A \tilde{G}^{A\mu\nu}. \quad (18)$$

Compact objects may also constrain parts of the dark-sector parameter space, although the resulting bounds are necessarily model dependent. Neutron stars can gravitationally focus DM from their environments [136–139], but focusing alone does not imply capture: An incoming dark particle must lose energy through scattering, dark self-interactions, or another dissipative process in order to become bound. In the present model, such interactions can arise from the portal operators coupling the Higgs field to dark glue in Eq. (18), as well as from the confined dark dynamics itself.

If the DM is made of bosonic glueballs and is captured efficiently, its behavior inside the star can differ substantially from that of ordinary nuclear matter. The characteristic glueball size is set by the confinement scale,

$$R_D \sim \Lambda_D^{-1} \simeq 2 \times 10^{-14} \text{ cm} \left( \frac{1 \text{ GeV}}{\Lambda_D} \right). \quad (19)$$

For  $\Lambda_D \sim \text{GeV}$ , the bound states are microscopic and carry internal momenta of order GeV, larger than the typical neutron Fermi momenta in neutron-star interiors. Moreover, glueballs are bosons and are not supported by Fermi pressure. Once captured and thermalized, a sufficiently large population may therefore settle near the stellar center and form a Bose condensate. Efficient annihilation of DM may heat up the neutron star, offering a potential

<sup>6</sup>In our model with vectorlike fermions,  $\bar{\theta}_D$  is physical, unlike in  $SU(2)_L$  of the SM, where due to the chiral nature of fermions, the electroweak  $\theta$  parameter can be rotated away [133–135].

probe in James Webb Space Telescope observations of old neutron stars [140–142].

A genuine neutron-star constraint requires the full sequence of capture, thermalization, condensation, and possible gravitational collapse and black-hole growth [140,143–147]. For a non-self-interacting boson of mass  $m_{\text{DM}}$ , the characteristic collapse number scales as  $N_{\text{crit}} \sim M_{\text{Pl}}^2/m_{\text{DM}}^2$ . Old neutron stars can therefore exclude part of the parameter space only if capture and thermalization are efficient and if any black hole formed grows faster than it evaporates. The relevant rates depend on the portal scale  $\Lambda$ , the confinement scale  $\Lambda_D$ , and the non-perturbative matrix elements controlling glueball interactions. In the minimal model,  $\Lambda$  is not fixed by the neutrino texture; it parametrizes whatever UV physics generates the higher-dimensional operators coupling the Higgs field to dark glue in Eq. (18). Thus, neutron-star bounds can be powerful, but they require a dedicated analysis of the portal strength, capture, thermalization, and collapse conditions. We leave this dark-sector phenomenology for future work.

Finally, the dark confinement transition occurs at a temperature of order  $T_D \sim \Lambda_D$ . For the minimal  $SU(2)_D$  theory, one should not assume that the transition is strongly first order. Pure  $SU(2)$  Yang-Mills theory has a second-order deconfinement transition in the three-dimensional Ising universality class [148,149], and light dark-sector fermions can further change the thermal history, depending on other details, such as the reheating temperature. Other theories, such as larger  $SU(N)_D$  groups or deformed dark sectors, may instead have first-order transitions. In such cases, the same dark gauge framework can produce a stochastic gravitational-wave background [150–152].

Precise predictions on the phase transition and the potential resulting gravitational waves require a finite-temperature analysis of the dark theory, which is left for future work.

*Conclusions.* We have proposed a minimal dark  $SU(2)$  model where the lightest neutrino mass is exactly zero and is protected by gauge symmetry. We derived the phenomenological predictions in the neutrino sector for this model and showed that it can be probed with future improvements in direct  $\beta$ -decay searches and cosmological neutrino mass measurements. Additionally, if a flavor symmetry fixes the right-handed neutrino mass basis, the neutrino mass matrix has a four-zero texture that could also be tested in ongoing and future oscillation experiments. Beyond the neutrino texture, the same anomaly-safe completion leads to a secluded confining sector with a potential DM candidate and, depending on the model parameters, may provide a promising target for dark-sector phenomenology using gravitational-wave and compact-object studies.

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*Data availability.* There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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