

Dynamical origin of neutrino masses and dark matter from a new confining sector

Maximilian Berbig^{1,2,*}, Juan Herrero-Garcia^{1,2,†} and Giacomo Landini^{2,‡}

¹*Departament de Física Teòrica, Universitat de València, 46100 Burjassot, Spain*

²*Instituto de Física Corpuscular (CSIC-Universitat de València), Parc Científic UV, C/Catedrático José Beltrán, 2, E-46980 Paterna, Spain*



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A dynamical mechanism, based on a confining non-Abelian dark symmetry, which generates Majorana masses for hyperchargeless fermions, is proposed. We apply it to the inverse seesaw scenario, which allows us to generate light neutrino masses from the interplay of TeV-scale pseudo-Dirac mass terms and a small explicit breaking of lepton number. A single generation of vectorlike dark quarks, transforming under a $SU(3)_D$ gauge symmetry, is coupled to a real singlet scalar, which serves as a portal between the dark quark condensate and three generations of heavy sterile neutrinos. Such a dark sector and the Standard Model (SM) are kept in thermal equilibrium with each other via sizable Yukawa couplings to the heavy neutrinos. In this framework, the lightest dark baryon, which has spin $3/2$ and is stabilized at the renormalizable level by an accidental dark baryon number symmetry, can account for the observed relic density via thermal freeze-out from annihilations into the lightest dark mesons. These mesons, in turn, decay to heavy neutrinos, which produce SM final states upon decay. This model may be probed by next generation neutrino telescopes via neutrino lines produced from dark matter annihilations.

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I. INTRODUCTION

The tiny observed mass scale of the active neutrinos is often explained via dynamical mechanisms that avoid the need to consider a small *ad hoc* Yukawa coupling by hand. Typically one either suppresses the contribution to m_ν by the mass scale of a heavy mediator or one introduces separate scalars with vacuum expectation values (VEVs) below the electroweak scale; see, for instance, Ref. [1]. In this work, we consider the inverse seesaw mechanism [2,3] (see also Ref. [4]), where one combines both aforementioned ingredients by invoking heavy vectorlike fermionic messengers, denoted as N from now on, and introduces a small explicit or spontaneous breaking μ of lepton number. Its main advantage, compared to the usual high-scale seesaw paradigm, is that it may be tested in low-scale experiments. However, note that the mechanism proposed in this work may be also used to generate the high-scale sterile neutrino Majorana mass present in the type I/III

seesaw [5–11], as well as light sterile Majorana masses (e.g., eV to keV scale).

Parametrically, in the inverse seesaw scheme the resulting active neutrino mass scales as

$$m_\nu \simeq 0.05 \text{ eV} \cdot y_\nu^2 \left(\frac{\mu}{1 \text{ keV}} \right) \left(\frac{35 \text{ TeV}}{M_D} \right)^2, \quad (1)$$

where M_D is the Dirac mass connecting both chiralities of N , and y_ν is the Yukawa coupling of N to the active neutrinos. While the original formulation of this mechanism in Ref. [2], which was based on additional singlet fermions added to a supersymmetric E_6 grand unified theory [12], assumed that this μ might arise from supersymmetry breaking, its relative smallness is often left unexplained. However, models with additional fermion singlets that generate μ radiatively have been proposed, see Refs. [13–15].

Here we pursue the idea that such a small breaking of lepton number might arise not from the VEV of an elementary scalar but *dynamically* from the formation of a dark quark condensate in a non-Abelian dark gauge theory. This idea was first applied in Ref. [16] to active neutrino masses in a field theory context by employing QCD and higher-dimensional operators connecting neutrinos to the quark condensate (see Ref. [17] for a string-theoretic realization and Refs. [18–21] for further investigations of the associated phenomenology). However such an approach requires at

*Contact author: berbig@ific.uv.es

†Contact author: juan.herrero@ific.uv.es

‡Contact author: giacomo.landini@ific.uv.es

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least one generation of massless quarks (unless additional model building steps are undertaken [21]), since the quarks are charged under the same symmetry that ensures the absence of the renormalizable neutrino mass terms, which is, however, heavily disfavored by recent lattice studies [22,23].

A compelling alternative is to consider an additional confining gauge group, which only acts on the dark sector, and to use this to generate a small Majorana mass for one chirality of the gauge singlet N , which evades the complications accompanying electroweak representations. Additionally the required dark sector can have the right ingredients for a successful dark matter (DM) candidate, which in our case is the lightest dark baryon \mathcal{B} . The presence of N with unsuppressed Yukawa couplings to both sectors turns out to be crucial for maintaining thermal equilibrium between them, which allows us to produce the dark matter abundance via thermal freeze-out of annihilations between \mathcal{B} and the lightest dark meson \mathcal{M} . References [24,25] used a similar but scale-invariant setup for generating electroweak symmetry breaking. When it comes to neutrino masses in scale-invariant frameworks, Ref. [26] generated (at most) TeV-scale right-handed neutrino masses required for realizations of the low scale [27,28] type I seesaw [5–8,10] from dynamical chiral symmetry breaking. The authors of Refs. [29,30] used a similar idea in order to explain a right-handed neutrino mass scale of $\mathcal{O}(10^7 \text{ GeV})$ in the context of the neutrino option [31,32]. Some of the aforementioned works considered dark matter in the form of dark pions produced via, e.g., the Higgs portal, whereas we consider dark baryons with the neutrino portal playing an important role in their thermalization with the Standard Model (SM) (for asymmetric DM connected to the SM via the neutrino portal, see Refs. [33–35]). Also our framework is not scale invariant and relies on a positive mass squared m_σ^2 for the singlet scalar.

This article is structured as follows: In Sec. II we introduce the field content, in Sec. III we specify the particle spectrum, and in Sec. IV we discuss the cosmological history. In Sec. V we elaborate on possible signatures for indirect detection and in Sec. VI we summarize our main results.

II. THE MODEL

In order to draw as much as possible from our knowledge of QCD, we choose the confining gauge group $\text{SU}(3)_D$ under which only the single vectorlike pair of dark quarks (q_L, q_R) transform in the fundamental representation. We further add three generations of vectorlike gauge singlet neutrinos (N_L, N_R) and the real¹ scalar σ . A \mathcal{Z}_4 symmetry

¹This simplification is just to avoid the presence of a Majoron [36,37] and not crucial to our mechanism.

TABLE I. Field content of the dark sector and SM leptons. We impose the gauged $\text{SU}(3)_D$ and discrete \mathcal{Z}_4 , whereas $\text{U}(1)_D$ is a residual symmetry present only after the scalar σ condenses. All dark fields are SM singlets.

	$\text{SU}(3)_D$	\mathcal{Z}_4	$\text{U}(1)_D$	Generations
q_L	3	$-i$	1	1
q_R	3	i	1	1
N_L	1	i	0	3
N_R	1	i	0	3
L	1	i	0	3
e_R	1	i	0	3
σ	1	-1	0	1

forbids bare Majorana masses for $N_L(N_R)$ as well as a bare mass term for the dark quarks (unlike the model in Ref. [38]). The real singlet scalar σ transforms as a -1 under \mathcal{Z}_4 and has a positive mass squared $m_\sigma^2 > 0$. The relevant terms read

$$-\mathcal{L}_{LN} = y_e \bar{L} H e_R + y_\nu \bar{L} \tilde{H} N_R + M_D \bar{N}_L N_R + \text{H.c.}, \quad (2)$$

$$-\mathcal{L}_D = y_Q \sigma \bar{q}_L q_R + y_{N_L} \sigma \bar{N}_L^c N_L + y_{N_R} \sigma \bar{N}_R^c N_R + \text{H.c.}, \quad (3)$$

$$V_\sigma = (m_\sigma^2 + \lambda_\sigma \sigma^2 + \lambda_{H\sigma} |H|^2) \sigma^2, \quad (4)$$

where we suppressed flavor indices and H is the SM Higgs doublet ($\tilde{H} \equiv i\sigma_2 H^*$) with a VEV v_H for its neutral component and all relevant charges and representations are summarized in Table I. Gauge confinement of $\text{SU}(3)_D$ generates a quark condensate

$$\langle \bar{q}_L q_R \rangle \simeq \Lambda_D^3, \quad (5)$$

where Λ_D is the confinement scale of the dark sector.² This induces a vacuum expectation value for the new scalar σ

$$\langle \sigma \rangle \simeq y_Q \frac{\Lambda_D^3}{m_\sigma^2}. \quad (6)$$

Such an ‘‘induced’’ VEV for σ is reminiscent of the well-known type II seesaw [39–45]. In our study we typically obtain $\langle \sigma \rangle \ll m_H \ll m_\sigma$, where m_H is the SM Higgs mass, so we can neglect mixing in the Higgs sector even for non-negligible $\lambda_{H\sigma}$. A nonvanishing $\langle \sigma \rangle$ implies Majorana

²RG running of the dark gauge coupling $\alpha_D = g_D^2/4\pi$ is given at one loop by

$$\frac{1}{\alpha_D(E_2)} = \frac{1}{\alpha_D(E_1)} + \frac{\beta_0}{2\pi} \log\left(\frac{E_2}{E_1}\right), \quad \beta_0 = \frac{11}{3} N_c - \frac{2}{3} = \frac{31}{3}.$$

The confinement scale is defined as the energy at which the dark gauge coupling turns nonperturbative, $\alpha_D(\Lambda_D) \simeq 4\pi$.

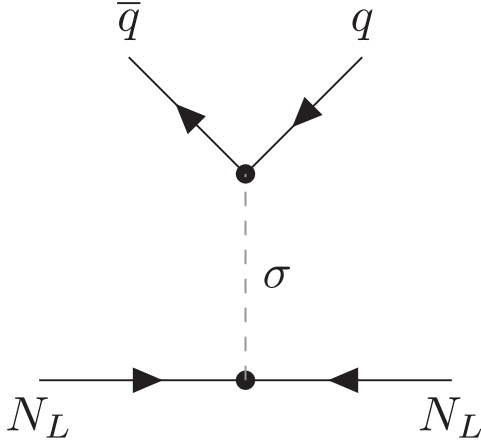


FIG. 1. Diagrammatic representation of the origin for the Majorana mass for $\overline{N}_L^c N_L$ from the chiral quark condensate $\langle \overline{q}q \rangle$. The same diagram also exists for $\overline{N}_R^c N_R$. The mechanism may be used to generate Majorana masses of heavy singlet/triplet fermions in seesaw type I/III, as well as the μ term in inverse seesaw.

masses for $N_{L,R}$ from the Feynman diagram shown in Fig. 1,

$$\mu_{L,R} \simeq 10 \text{ keV} \cdot y_Q y_{N_{L,R}} \left(\frac{\Lambda_D}{10 \text{ TeV}} \right)^3 \left(\frac{3.1 \times 10^8 \text{ GeV}}{m_\sigma} \right)^2. \quad (7)$$

Note that the inverse seesaw requires only one of the two mass terms $\mu_{L,R}$ to be present to induce a small mass for m_ν ; if both terms are nonzero, one obtains that at leading order in the seesaw expansion $\mu_{L,R} \ll y_\nu v_H \ll M_D$ [46]

$$m_\nu \simeq \frac{\mu_L y_\nu^2 v_H^2}{M_D^2 - \mu_L \mu_R}. \quad (8)$$

For the remainder of this work we define $\mu \equiv \mu_L$ and ignore the subleading correction in the denominator, hence we use $m_\nu \simeq \mu (y_\nu v_H / M_D)^2$. The small Majorana masses $\mu_{L,R}$ are responsible for splitting the masses of the two chiralities of N so that they become ‘‘pseudo-Dirac’’ fermions.

The spontaneous breaking of the \mathcal{Z}_4 symmetry leads to the formation of domain walls. To avoid that, we break this symmetry explicitly in the scalar potential by the bias term $\kappa \sigma |H|^2$ [47]. We assume a negligible trilinear term $\mu_3 \sigma^3$ and the linear piece $\mu_1^3 \sigma$ can always be set to zero by a field redefinition [48]. The dimensionful coupling κ does not induce a sizable shift to $\langle \sigma \rangle$ in Eq. (6) as long as

$$\kappa \ll 3 \times 10^7 \text{ GeV} \cdot y_Q \left(\frac{\Lambda_D}{10 \text{ TeV}} \right)^3, \quad (9)$$

and the domain walls decay before big bang nucleosynthesis (BBN) for

$$\kappa \gtrsim 10^{-13} \text{ GeV} \cdot y_Q \left(\frac{\Lambda_D}{10 \text{ TeV}} \right)^3 \left(\frac{10^8 \text{ GeV}}{m_\sigma} \right). \quad (10)$$

We checked that decays before BBN occur long before the domain walls dominate the energy budget of the Universe. Another way to remove the domain walls could be to invoke the fact that the \mathcal{Z}_4 symmetry is anomalous with respect to $\text{SU}(3)_D$ [49], but that requires a dedicated study.

III. PARTICLE SPECTRUM

At temperatures below Λ_D the theory confines and the dark quarks reorganize into dark hadrons. Since, unlike in real QCD, we work with only one generation of dark quarks, there are no spontaneously broken chiral symmetries and hence no Nambu-Goldstone modes similar to the pions [50]. The only global symmetries would be

$$\text{U}(1)_D \otimes \text{U}(1)_A, \quad (11)$$

where the vectorial $\text{U}(1)_D$ is the dark equivalent of baryon number and the axial symmetry $\text{U}(1)_A$ is both explicitly broken by the coupling to σ and anomalous to begin with. We denote the lightest meson state $|\overline{q}q\rangle$ as \mathcal{M} , which is parity odd [51]. Since the coupling to σ explicitly breaks a chiral symmetry, one expects a dark quark mass of $m_Q \simeq y_Q \langle \sigma \rangle$ [26]. Because of the absence of spontaneously broken chiral symmetries and thus dark pions, the equivalent of the Gell-Mann–Oakes–Renner relation [52] $m_\pi \sim \sqrt{m_Q \Lambda_D}$ is not valid for the meson \mathcal{M} and to reduce the number of free parameters we fix

$$m_{\mathcal{M}} \simeq \Lambda_D. \quad (12)$$

If we added more than one flavor of dark quarks, we would find that the resulting lighter dark pions would typically be so long-lived that the injection of electromagnetic radiation from the electroweak showers of their leptonic decay products would alter BBN significantly [53] or they could be stable enough to overclose the Universe.³

The lightest dark baryon \mathcal{B} consisting of $|qqq\rangle$ has spin $3/2$ [54,55], similar to the Δ resonance of the strong interaction, and its mass is expected from large- N_c (number of dark colors) arguments to scale as $N_c \Lambda_D$ [54]. In analogy to QCD, Ref. [55] finds a mass scale of about $10 \Lambda_D$ and we interpolate between these two estimates by setting

³From the Gell-Mann–Oakes–Renner relation the dark pions would get a tiny mass $m_\pi \simeq y_Q \Lambda_D^2 / m_\sigma$ typically below M_D . Thus they would have to decay into active neutrinos with a highly suppressed rate $\sim m_\nu^2 \Lambda_D^4 / m_\sigma^5$. A possible way to generalize our model to $N_f \geq 2$ generations of dark quarks is to introduce a bare quark mass term, which softly breaks the \mathcal{Z}_4 symmetry. This generates a larger mass for the dark pions above M_D allowing them to decay into N before BBN.

$$m_B \simeq 5\Lambda_D. \quad (13)$$

The stability of \mathcal{B} at the renormalizable level is ensured by the conservation of dark baryon number $U(1)_D$ and hence it may be a good DM candidate.

For $SU(3)_D$ there can be glueball states that are odd or even under the dark sector charge conjugation. The even glueballs can decay into two mesons as long as their mass is above $2m_M$ in analogy to what is expected for QCD glueballs [56,57]. Odd glueballs might be stable and could be produced by similar dynamics to the dark baryons (see Sec. IV), forming a component of DM [58–60]. In the absence of more detailed knowledge of the mass spectrum for the case at hand, which would certainly require a lattice simulation, we focus on the dark baryon as DM and assume that the odd glueballs decay away. Even if these glueballs were stable, they would correct our results by only $\mathcal{O}(1)$ factors.⁴ If we chose the adjoint of $SU(3)_D$ for the representation of the dark quarks, there could be exotic hybrid bound states of dark quarks and dark gluons playing the role of DM [61].

IV. COSMOLOGICAL HISTORY

We set both the reheating temperature T_{RH} and the maximum temperature during reheating T_{max} , which can be much larger than T_{RH} [62,63] for noninstantaneous reheating, to be smaller than m_σ so that we can safely integrate σ out and treat it purely as a mediator. Reheating sets the stage for a thermalized SM plasma and we assume that $T_{RH} \gg M_D, \Lambda_D$. As long as the N are relativistic, they are produced via their sizable Yukawa coupling to the SM leptons at temperatures below

$$T_N^{in} \simeq 7 \times 10^7 \text{ GeV} \cdot \left(\frac{y_\nu}{10^{-4}} \right)^2 \sqrt{\frac{106.75}{g_{*\rho}(T_N^{in})}}, \quad (14)$$

where $g_{*\rho}$ is the number of relativistic degrees of freedom in the energy density and they consequently populate the dark sector via their fast σ -mediated annihilations into dark quarks $NN \leftrightarrow \bar{q}q$ at $T > \Lambda_D, M_D$. Dark gluons are produced from the thermalized quarks via $\bar{q}q \leftrightarrow gg$ with a rate of about $g_D^4 T$. Requiring that this comes into thermal equilibrium above the dark confinement scale amounts to the requirement

$$\Lambda_D \lesssim 10^{10} \text{ GeV} \cdot \left(\frac{g_D}{0.01} \right)^4 \sqrt{\frac{106.75}{g_{*\rho}(T_{gg}^{in})}}, \quad (15)$$

⁴Glueballs have a typical mass $m_{DG} \simeq \Lambda_D$ and annihilate with cross section $\sigma_{DG} \simeq 1/\Lambda_D^2$, leading to qualitatively the same dynamics as for the dark baryons.

which, as we will see in Eq. (18), is always satisfied, especially for nonperturbative g_D .

The abundance of dark baryons is determined via thermal freeze-out of the annihilations $\bar{\mathcal{B}}\mathcal{B} \leftrightarrow \bar{\mathcal{M}}\mathcal{M}$, which occurs in the s wave⁵ with a geometric cross section [58] of [65]

$$\langle \sigma_D |\bar{v}| \rangle \simeq \frac{\pi}{\Lambda_D^2} \simeq 2.2 \times 10^{-26} \frac{\text{cm}^3}{\text{s}} \cdot \left(\frac{41 \text{ TeV}}{\Lambda_D} \right)^2. \quad (16)$$

Here we neglect any potential enhancement of the cross section due to intermediate resonances, as occurs, e.g., for the proton due to the deuteron resonance [66]. To logarithmic accuracy we find that these annihilations decouple at

$$T_{BM}^{\text{out}} \simeq \frac{m_B}{25} \quad (17)$$

and that the relic abundance, in general, is reproduced for values of [54,55]

$$\Lambda_D \simeq (1-100) \text{ TeV}. \quad (18)$$

We stress that, due to $\mathcal{O}(1)$ uncertainties in the hadron spectrum, as well as in the thermally averaged cross section of Eq. (16), we are unable to determine the precise value of Λ_D that reproduces the relic abundance and can only estimate a reasonable range in Eq. (18). Even though the dark baryon is self-interacting with $\sigma_B/m_B \sim 1/\Lambda_D^3$, the above range for Λ_D precludes a strong enough elastic cross section that would be necessary to solve the ‘‘cuspcore’’ [67–70] or ‘‘too-big-to-fail’’ problems [71,72] (for an overview, see Ref. [66] and references within), which would require sub-GeV values of Λ_D . For the same reason, our scenario is not constrained by bounds from the halo ellipticity [73] or the Bullet Cluster [74].

A crucial ingredient of the freeze-out estimate is that the dark sector maintains the same temperature as the SM bath. Indeed, below the confinement scale, all the dark hadrons are nonrelativistic. Therefore, in absence of thermal contact with the SM, the temperature of the dark thermal bath would redshift only logarithmically with the scale factor, analogously to ‘‘cannibal’’ DM models [75,76], leading to an overabundant population of dark baryons. Thermal

⁵Following the argument in Refs. [61,64], the maximum angular momentum of the annihilation process can be estimated as $l = \mu |\bar{v}| b$, where $b \simeq 1/\Lambda_D$ is the impact parameter, $|\bar{v}| \simeq \sqrt{T/m_B}$ the relative velocity, and $\mu = m_B/2$ is the reduced mass. As $l \simeq \sqrt{T m_B}/(2\Lambda_D)$ we find $l < 1$ both at the time of DM freeze-out [$l \simeq 1/2$ see Eq. (17)] and today, so the relevant processes always take place in the fully quantum regime. In this regime, the s -wave contribution is expected to dominate the total annihilation cross section $\langle \sigma_D |\bar{v}| \rangle$ [61]. Notice that in [61,64] some of the processes take place in the semiclassical regime because they consider hadrons made of heavy quarks.

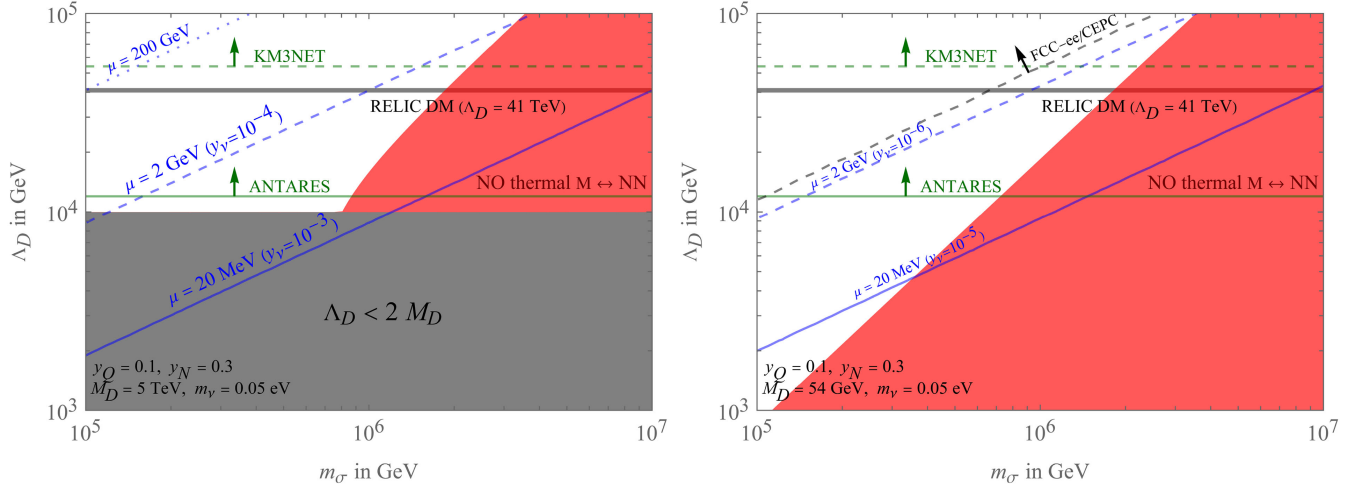


FIG. 2. Allowed parameter space for the successful production of the dark matter together with isocontours for μ and the corresponding y_ν required to explain $m_\nu = 0.05$ eV in the inverse seesaw mechanism for two sets of benchmark parameters with TeV-scale (left) and GeV-scale (right) N . The green lines indicate the current and projected lower limits on Λ_D from indirect detection via dark matter annihilations into dark mesons, followed by a chain of two-body decays producing neutrinos. For N around the GeV-scale, the parameter space might also be probed by displaced vertex searches at future colliders.

equilibrium (between both sectors⁶) below the confinement scale is maintained via decays and inverse decays $\mathcal{M} \leftrightarrow \bar{N}N$, because the N are tightly coupled with the SM at temperatures above their mass M_D , see Eq. (14). One finds a decay rate of

$$\Gamma(\mathcal{M} \rightarrow \bar{N}N) \simeq \frac{y_D^2 y_N^2 m_{\mathcal{M}} f_{\mathcal{M}}^2}{32\pi m_\sigma^4} \sqrt{1 - \frac{4M_D^2}{m_{\mathcal{M}}^2}}, \quad (19)$$

in terms of the matrix element $f_{\mathcal{M}}$, that we parametrize as

$$f_{\mathcal{M}} \equiv \langle 0 | \bar{q} \gamma_5 q | \mathcal{M} \rangle \simeq \Lambda_D^2, \quad (20)$$

and $y_N \equiv y_{N_L} + y_{N_R}$. By employing Maxwell-Boltzmann statistics, we find that the thermally averaged decay rate for $\mathcal{M} \rightarrow \bar{N}N$ reads [78]

$$\langle \Gamma_D \rangle = \frac{K_1\left(\frac{m_{\mathcal{M}}}{T}\right)}{K_2\left(\frac{m_{\mathcal{M}}}{T}\right)} \Gamma(\mathcal{M} \rightarrow \bar{N}N) \quad (21)$$

in terms of the modified Bessel functions of the second kind $K_{1,2}(m_{\mathcal{M}}/T)$. From the principle of detailed balance, we obtain for the thermal average of the inverse decay rate $\bar{N}N \rightarrow \mathcal{M}$,

⁶An alternative possibility is that the dark sector particles thermalize among themselves, but not with the SM, forming a secluded dark bath that evolves with its own temperature $T_d \equiv \xi T$ [77]. In such a case the DM evolution depends on the additional free parameter ξ , which encodes some unknown initial condition, so that $\xi \simeq 10^{-3}$ is needed to reproduce the correct relic abundance for TeV-scale DM. We do not consider this possibility here.

$$\langle \Gamma_{\text{ID}} \rangle = \frac{n_{\mathcal{M}}^{\text{eq}}}{n_N^{\text{eq}}} \langle \Gamma_D \rangle. \quad (22)$$

Here we introduced the equilibrium number density of particles species i with mass m_i and g_i internal degrees of freedom via

$$n_i^{\text{eq}} = \frac{g_i T^3}{2\pi^2} \left(\frac{m_i}{T}\right)^2 K_2\left(\frac{m_i}{T}\right), \quad (23)$$

with $g_{\mathcal{M}} = 1$, $g_N = 2$. We require that the inverse decay remains in equilibrium until at least the temperature of the \mathcal{B} freeze-out defined in Eq. (17). On top of that we impose $m_{\mathcal{M}} = \Lambda_D > 2M_D$ so that the decay channel is kinematically open. Further, we check that the dark matter relic density is not diluted by entropy release: we find that decay $\mathcal{M} \rightarrow \bar{N}N$ is always fast, as expected from the previous arguments, and that the decay width of N given by $\Gamma(N \rightarrow LH) \simeq y_N^2 M_D / (8\pi)$ equals the Hubble rate before the temperature $T_{\text{dom}} \simeq 7M_D / (4g_{*p}(T_{\text{dom}}))$ [79] (when N would start to dominate over the energy density of radiation) as long as

$$y_\nu \gtrsim 7.8 \times 10^{-9} \cdot \sqrt{\frac{M_D}{10 \text{ TeV}}} \left(\frac{g_{*p}(T_{\text{dec}})}{106.75}\right)^{\frac{1}{4}} \left(\frac{106.75}{g_{*p}(T_{\text{dom}})}\right). \quad (24)$$

In Fig. 2 we depict the parameter space in the Λ_D versus m_σ plane subject to the previously discussed constraints. For the showcased benchmark point, we obtain the observed relic density of $\Omega_B h^2 = 0.120 \pm 0.001$ [80] by using Eq. (16) and values of $\mu > \mathcal{O}(100 \text{ MeV})$ corresponding

to $y_\nu < \mathcal{O}(10^{-3})$, bounded from below by Eq. (24). These values for the lepton number breaking parameter are larger than the conventionally assumed keV scale [see, e.g., Eq. (1)], but bear the additional advantage that $M_D = \mathcal{O}(\text{TeV})$, which is in range of future collider experiments. One can understand the largeness of μ from the plot in Fig. 2 by noting that for smaller values of μ (equivalent to larger y_ν for fixed m_ν) the relic abundance would only be reached in the red colored region where \mathcal{M} is not thermalized long enough. For values of μ above the scale of M_D , the inverse seesaw expansion breaks down and we would be in the usual type I seesaw regime. A smaller value for the product $y_Q y_N$ moves the μ isocontours upward along the Λ_D axis and increases the size of the region excluded by the meson thermalization. Finally, let us point out that our scheme involving TeV-scale N might reproduce the baryon asymmetry of the Universe [81] via resonantly enhanced [82–84] out-of-equilibrium decays.

V. SIGNALS AND CONSTRAINTS

For our parameter space with $\mu > \mathcal{O}(10 \text{ MeV})$, the mixing between active neutrinos of flavor $i = e, \mu, \tau$, and N scales as $|V_{iN}|^2 \simeq m_\nu/\mu \lesssim \mathcal{O}(10^{-[9,8]})$. A recent review of all pertinent laboratory constraints was compiled in Ref. [85]: For $M_D = \mathcal{O}(\text{TeV})$ the constraint from electroweak precision observables due to the modification of charged- and neutral-current reactions induced by the

nonunitarity in the active neutrino sector reads $|V_{iN}|^2 < 10^{-3}$ [86–88]. Next generation electron colliders such as FCC-ee or CEPC could improve the nonunitarity bound for the mixing with electron neutrinos down to $|V_{eN}|^2 < 10^{-[5,4]}$ [89]. Displaced vertex searches for N masses below about 100 GeV at the high luminosity upgrade of the LHC might probe values down to $|V_{iN}|^2 \simeq 5 \times 10^{-10}$, while the proposed hadronic collider FCC-hh might reach $|V_{iN}|^2 \simeq 5 \times 10^{-11}$ [89], whereas FCC-ee or CEPC could potentially test mixings as small as $|V_{iN}|^2 \simeq 10^{-11}$ [90–92]. These searches for long-lived N far below the TeV scale could test our cosmologically preferred parameter space (see the right plot in Fig. 2), as can be observed in Fig. 3.

When it comes to charged lepton flavor violation, the strongest constraints come from the nonobservation of the decay $\mu \rightarrow e\gamma$ at MEG [93] and MEG II [94], setting a combined limit of 3.1×10^{-13} [95], and nonobservation of muon-to-electron conversion on titanium by the SINDRUM II Collaboration with an upper limit on the branching ratio for ground state transitions of $R_{\mu \rightarrow e}^{\text{Ti}} < 4.3 \times 10^{-12}$ [96], that impose $|V_{eN} V_{\mu N}^*| < 10^{-5}$ in the range $100 \text{ GeV} < M_D < 10 \text{ TeV}$ [85]. At two loops in the inverse seesaw [97,98], there are contributions to the electric dipole moment of the electron, which is measured by the ACME II experiment to be $|d_e| < 1.1 \times 10^{-29} e \text{ cm}$ [99] and in the future this result is expected to improve by about an order of magnitude. The authors of Ref. [100] found for $M_D = \mathcal{O}(\text{TeV})$ the largest

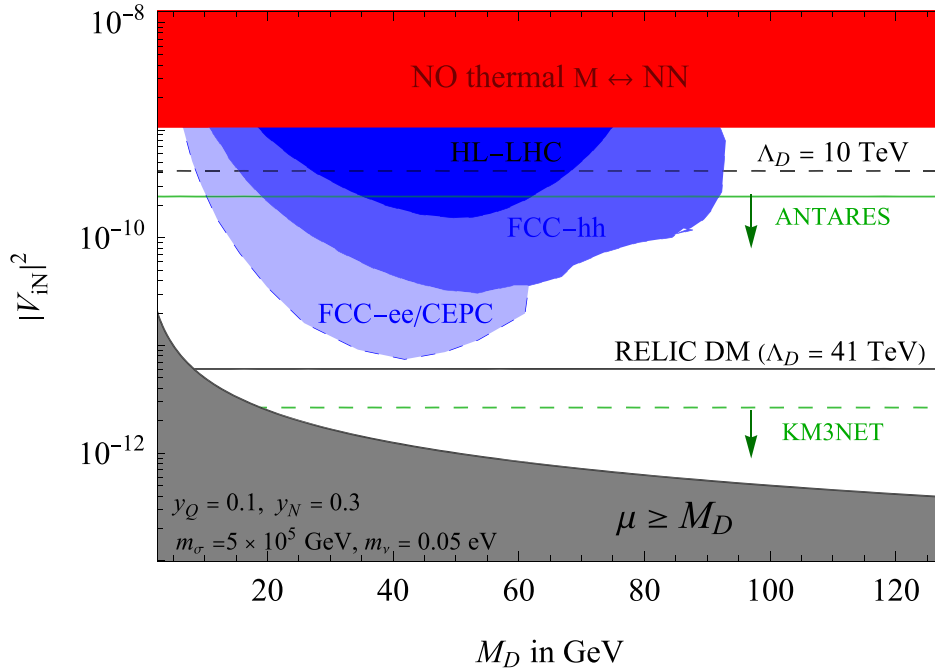


FIG. 3. Sensitivity projections of displaced vertex searches for GeV-scale N together with limits and projections from neutrino telescopes in the parameter space that is compatible with our cosmological considerations. Our cosmological analysis is independent of flavor and the collider limits are for the electron and muon channels. In the gray region, the inverse seesaw scheme breaks down as the Majorana masses μ of N become larger than their Dirac masses M_D .

possible value of $|d_e| = 10^{-[32,31]}$ e cm, which could be testable in the future. For the muon electric dipole moment, the current direct limit of $|d_\mu| < 1.9 \times 10^{-19}$ e cm [101] was obtained at Brookhaven National Laboratory (BNL) and indirect limits from heavy atoms and molecules via muon loops reach down as far as 2×10^{-20} e cm [102], whereas future experiments at PSI and J-PARC are expected to improve these bounds to 6×10^{-23} [103–105] and 10^{-24} e cm [106], respectively. By rescaling the result from Ref. [100] for the electron electric dipole moment in the inverse seesaw by a factor of m_μ/m_e we find $|d_\mu| \simeq 2 \times 10^{-[30,29]}$ e cm, which is out of reach of future experiments. In Ref. [107] it was concluded that the inverse seesaw cannot account for the discrepancy in the anomalous magnetic dipole moment of the muon observed by BNL [108].

For N above the GeV scale one expects that the dominant contribution to neutrinoless-double- β decay comes from the exchange of the light neutrinos. An estimate of the resulting rate depends on assumptions about flavor, and for the normal hierarchy of active neutrino masses there remains the possibility of an accidental cancellation drastically reducing the rate below future sensitivities, depending on the interplay of so far unknown lightest neutrino mass, the observed mixing angles, and the possible Majorana phases in the Pontecorvo-Maki-Nakagawa-Sakata matrix [109].

There will not be a signal in the gravitational wave spectrum from the minimal dark confinement transition considered here, since for three dark colors and one dark flavor one expects a smooth crossover (see Ref. [110] and references within). Our setup is not constrained by direct detection of DM due to the smallness of the mixing between σ and the SM Higgs boson [see the discussion below Eq. (6)].

Indirect detection on the other hand, offers the intriguing prospect of signals at future neutrino telescopes: Annihilations of \mathcal{B} into mesons followed by the immediate decay of \mathcal{M} into two TeV-scale N , which in turn is followed by the decays $N \rightarrow \sum_i H_0 \nu_i, \sum_i Z_\mu \nu_i, \sum_i W_\mu^\mp l_i^\pm$ with $\text{BR}(N \rightarrow \sum_i H_0 \nu_i) = \text{BR}(N \rightarrow \sum_i Z_\mu \nu_i) = \text{BR}(N \rightarrow \sum_i W_\mu^\mp l_i^\pm)/2$ [111], would produce monoenergetic primary neutrinos with an energy of $m_B/4$ (assuming $m_B \gg m_{\mathcal{M}}$ for simplicity) plus less energetic secondary neutrinos and a background of SM particles from the decays and interactions of the H_0, Z_0, W^\mp, l^\pm also present.⁷ For reactions producing initially monoenergetic neutrinos with energies above $\mathcal{O}(100 \text{ GeV})$, there would never be an exactly monochromatic neutrino line as the produced neutrinos would be so energetic that they would produce electroweak bremsstrahlung resulting in cascades similar to QCD jets at

colliders [112–114]. Reference [115] concluded that the widening of the line is not larger than the energy resolution of high energy neutrino telescopes [116] for dark matter masses in the range $10^{[3,8]}$ GeV, which is why in the following we neglect this effect.

Current and projected bounds on neutrinos produced in dark matter annihilations were compiled in Ref. [117] and for the $m_B = \mathcal{O}(10^{[4,5]})$ TeV mass range the strongest bound (for s -wave annihilations) for neutrino telescopes of $\langle \sigma |\vec{v}| \rangle < (5 \times 10^{-24} - 10^{-23}) \text{ cm}^3/\text{s}$ was obtained by the ANTARES Collaboration [118]. This limit is expected to improve to $10^{-24} \text{ cm}^3/\text{s}$ [119] for the expansion of the current IceCube observatory [120], $(10^{-25} - 5 \times 10^{-25}) \text{ cm}^3/\text{s}$ for the proposed Pacific Ocean Neutrino experiment [121], and $5 \times 10^{-[26,25]} \text{ cm}^3/\text{s}$ for the KM3NeT [122,123] water Cherenkov detector currently being constructed with a km^3 volume in the Mediterranean sea. From the aforementioned electroweak showers and the charged SM particles that can also be emitted in the N decays, we also expect that high energy photons should be produced, which is constrained by γ ray data from Fermi-LAT [124] and H.E.S.S. [125] leading to $\langle \sigma |\vec{v}| \rangle < 10^{-23} \text{ cm}^3/\text{s}$ [126,127] and will be probed further by the upcoming Cherenkov Telescope Array [128] with a projected limit of $(10^{-24} - 5 \times 10^{-24}) \text{ cm}^3/\text{s}$ [126,127]. Note that a detailed limit for our case will depend on the energy fraction deposited in the photons, which requires a dedicated simulation of the decay chain and showering [127].

All of the above limits assume spin 1/2 Majorana dark matter (2 degrees of freedom), but since we have a spin 3/2 Dirac fermion [$2 \times (2 \times 3/2 + 1) = 8$ degrees of freedom] we need to rescale the limits by a factor of 1/4. Also limits from annihilations assume a neutrino energy equal to the dark matter mass, but for our case the right energy range is roughly $m_B/4$. We estimate the thermally averaged dark matter annihilation cross section by using (16) and with the rescaled limit from ANTARES [118] we find

$$\Lambda_D > 12 \text{ TeV} \cdot \sqrt{\frac{0.25 \times 10^{-24} \text{ cm}^3/\text{s}}{\langle \sigma |\vec{v}| \rangle}}, \quad (25)$$

which will improve to 54 TeV once KM3NeT [122,123] is operational so that our benchmark of $\Lambda_D = 41 \text{ TeV}$ should be tested by next generation experiments; we demonstrate the impact on our parameter space in Fig. 2. Of course, we should stress that this is just an order of magnitude estimate due the uncertainties related to the nonperturbative dynamics of the dark hadrons and setting $\text{BR}(\mathcal{M} \rightarrow \bar{N}N) \simeq 1$.

If DM is stabilized by a global symmetry such as our accidental $U(1)_D$, there might be higher-dimensional operators from the putative field theoretic UV completion that break this symmetry or Planck-suppressed operators due to nonperturbative quantum gravitational effects, which are expected to violate all global symmetries [129–132] and

⁷For the case of light N with masses below the electroweak scale, the available decay modes of N depend on its mass, which will affect the number of produced neutrinos and their energy spectra. We do not expect the order of magnitude estimates obtained for TeV-scale N to drastically change in this limit.

thus induce DM decay [133]. We summarized the relevant effective operators for our case in the Appendix and find that the lowest-dimensional ones occur at dimension $d = 8$ in Eqs. (A10) and (A11). On dimensional grounds, we parametrize

$$f_B \equiv \langle 0|qqq|\mathcal{B}\rangle \simeq \Lambda_D^3. \quad (26)$$

The operator in Eq. (A10) induces the decays $\mathcal{B}_\mu \rightarrow \sum_i H_0 \nu_i$, $\mathcal{B}_\mu \rightarrow \sum_i Z_\mu \nu_i$ and $\mathcal{B}_\mu \rightarrow \sum_i W_\mu^\mp l_i^\pm$. Owing to the fact that $m_B \gg m_H, m_Z, m_W$ one finds that [134] $\Gamma(\mathcal{B}_\mu \rightarrow \sum_i H_0 \nu_i) = \Gamma(\mathcal{B}_\mu \rightarrow \sum_i Z_\mu \nu_i) = \Gamma(\mathcal{B}_\mu \rightarrow \sum_i W_\mu^\mp l_i^\pm)/2$ and their sum reads [134]

$$\Gamma_2^{(1)} = \frac{|c_8^{(1)}|^2 \Lambda_D^6 m_B^3}{\Lambda_{UV}^8 256\pi}. \quad (27)$$

There also exist the three-body modes $\mathcal{B}_\mu \rightarrow \sum_i H_0 Z_\mu \nu_i$ and $\mathcal{B}_\mu \rightarrow \sum_i H_0 W_\mu l_i$ that are enhanced by a factor of m_B^2/v_H^2 compared to the previous two-body decays and thus dominate over them. We estimate their sum to be

$$\Gamma_3^{(1)} \simeq \frac{|c_8^{(1)}|^2 \Lambda_D^6}{\Lambda_{UV}^8} \frac{3m_B^5}{8192\pi^2 v_H^2}. \quad (28)$$

In Ref. [115] it was found that the energy spectrum for a three-body decay can be approximated by a power law with $dN/dE \sim (E/m_B)^{-[2,3]}$ for which they derive a limit of about $\tau_B > 10^{28}$ s in the window $m_B = \mathcal{O}(10^4\text{--}10^5 \text{ GeV})$ using data from IceCube [135] and we obtain

$$\frac{\Lambda_{UV}}{10^{12} \text{ GeV}} \gtrsim |c_8^{(1)}|^{\frac{1}{4}} \left(\frac{m_B}{5\Lambda_D}\right)^{\frac{5}{8}} \left(\frac{\Lambda_D}{40 \text{ TeV}}\right)^{\frac{11}{8}} \left(\frac{10^{28} \text{ s}}{\tau_B}\right)^{\frac{1}{8}}. \quad (29)$$

The second operator in Eq. (A11) leads to the decay modes $\mathcal{B}_\mu \rightarrow A_\mu N$ and $\mathcal{B}_\mu \rightarrow Z_\mu N$ followed by the aforementioned two-body decay of N to SM states. The total width for our case of $m_B \gg M_D$ is found to be [134]

$$\Gamma_2^{(2)} = \frac{|c_8^{(2)}|^2 \Lambda_D^6 m_B^3}{\Lambda_{UV}^8 4\pi}. \quad (30)$$

Consequently, we expect monochromatic neutrino lines from the DM decays due the operator in Eq. (A11), while the previously mentioned, but negligible, widening of the lines due to electroweak cascades applies again. The authors of [115] derived a limit on the DM lifetime for two-body decays of about $\tau_B > 10^{28}$ s in the aforementioned window of DM masses. Applying this limit to the rate in (30) leads to the constraint

$$\frac{\Lambda_{UV}}{6 \times 10^{11} \text{ GeV}} > |c_8^{(2)}|^{\frac{1}{4}} \left(\frac{m_B}{5\Lambda_D}\right)^{\frac{3}{8}} \left(\frac{\Lambda_D}{40 \text{ TeV}}\right)^{\frac{9}{8}} \left(\frac{10^{28} \text{ s}}{\tau_B}\right)^{\frac{1}{8}}, \quad (31)$$

which is slightly weaker than Eq. (29), and the bound for the rate in Eq. (27) would be found by replacing $c_8^{(2)}$ with $c_8^{(1)}/8$.

VI. CONCLUSIONS

We have proposed a dark sector that dynamically generates the lepton number breaking mass term of an electrically neutral fermion via the condensation of a single generation of dark quarks. This may be used to generate either the Majorana mass of the heavy right-handed neutrinos present in type I/III seesaw⁸ or the small μ term present in the inverse seesaw mechanism. We focus on the latter possibility because it has a richer phenomenology.

The resulting massive spin 3/2 dark baryon which emerges from the confinement of the dark sector is analogous to the Δ baryon of QCD and is stabilized by a dark baryon number symmetry, which is only violated by higher-dimensional operators starting at dimension 8 due to its larger spin. Thus it constitutes a good DM candidate and we obtain its yield from thermal freeze-out in the dark sector. This model predicts a dark confinement scale of the order of $\Lambda_D = \mathcal{O}(1\text{--}100)$ TeV to reproduce the dark matter relic abundance from which we find that $\mu > \mathcal{O}(10 \text{ MeV})$ and a heavy neutrino mass of $M_D = \mathcal{O}(\text{TeV})$. The neutrino portal from the inverse seesaw is crucial for keeping the dark hadrons in thermal contact with the SM. We have studied the phenomenological implications of the TeV-scale pseudo-Dirac neutrinos and the potential signals from DM decays and annihilations, the latter of which may be probed by the upcoming KM3NeT experiment. Furthermore, a detailed analysis of the energy spectra of the decay products relevant for indirect detection would be interesting to pursue.

In the case of GeV-scale sterile neutrinos, displaced vertex searches at proposed future colliders FCC-ee and FCC-hh may probe a significant part of the allowed parameter space.

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⁸In principle, any Majorana mass scale (say, e.g., eV to 10^{14} GeV) could be realized. However, in order to generate very heavy Majorana masses, larger than $\mathcal{O}(100)$ TeV, and reproduce the relic abundance via freeze-out [which requires $\Lambda_D \lesssim \mathcal{O}(100)$ TeV] one needs $m_\sigma \ll \Lambda_D$, which is a different region of the parameter space to the one analyzed in this work.

APPENDIX: SPIN 3/2 FERMIONS AND DARK MATTER STABILITY

Here we briefly review the properties of spin 3/2 Rarita-Schwinger fields [136] in order to construct the interactions of our DM candidate $\mathcal{B} \sim |qqq\rangle$ which is analogous to the Δ baryon of QCD. We begin with a ‘‘spinor-vector’’ ψ^μ , which is the direct product of the vector representation $(\frac{1}{2}, \frac{1}{2})$ and the spinor representation of a Dirac fermion $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$, which gives

$$\left(1, \frac{1}{2}\right) \oplus \left(0, \frac{1}{2}\right) \oplus \left(\frac{1}{2}, 1\right) \oplus \left(\frac{1}{2}, 0\right), \quad (\text{A1})$$

and corresponds to 16 degrees of freedom (d.o.f.). By imposing the constraint on the free theory [137] that

$$\gamma_\mu \psi^\mu = 0, \quad (\text{A2})$$

one can eliminate the spin 1/2 Dirac spinor $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$, which corresponds to four d.o.f. Further imposing a second constraint [137],

$$\partial_\mu \psi^\mu = 0, \quad (\text{A3})$$

eliminates four more d.o.f. from another spin 1/2 Dirac spinor and the remaining eight physical d.o.f. correspond to a spin 3/2 Dirac fermion of mass m that obeys the Dirac equation. It has a kinetic term given by

$$\bar{\psi}^\mu \Lambda_{\mu\nu} \psi^\nu \quad (\text{A4})$$

in terms of [138]

$$\begin{aligned} \Lambda_{\mu\nu} = & -(\not{p} - m)g_{\mu\nu} + A(\gamma_\mu \not{p}_\nu + p_\mu \gamma_\nu) \\ & + \frac{1}{2}(1 + 2A + 3A^2)\gamma_\mu \not{p}_\nu + m(1 + 3A + 3A^2)\gamma_\mu \gamma_\nu. \end{aligned} \quad (\text{A5})$$

Here A is a free parameter with the requirement $A \neq -1/2$ to avoid a singular propagator [138] and the structure of the kinetic term can be obtained by requiring invariance under the following field redefinition [139,140]:

$$\psi^\mu \rightarrow (g^{\mu\nu} + a\gamma^\mu \gamma^\nu)\psi_\nu, \quad (\text{A6})$$

$$A \rightarrow \frac{A - 2a}{1 + 4a}, \quad (\text{A7})$$

in terms of another free parameter $a \neq -1/4$. One can think of A as parametrizing the admixture of the spin 1/2 component $\gamma_\mu \psi^\mu$ in the off shell ψ^μ field [139]. Because of the invariance of the Lagrangian under the above transformations the parameter A will drop out of all physical observables as shown by Ref. [141]. By imposing invariance under the aforementioned field redefinitions, one can construct the interaction of ψ^μ with fermions and

pseudoscalars, e.g., the coupling of Δ to pions π and nucleons n [139,140],

$$\mathcal{L}_{\Delta\pi n} = c_{\Delta\pi n} \Delta^\mu \theta_{\mu\nu} (\partial^\nu \pi) n + \text{H.c.}, \quad (\text{A8})$$

where we suppressed isospin and the electric charges and one defines [139,140]

$$\theta_{\mu\nu} \equiv g_{\mu\nu} + \left(\frac{A}{2}(1 + 4z) + z\right) \gamma_\mu \gamma_\nu, \quad (\text{A9})$$

where the free parameter z is known as the ‘‘off shell parameter,’’ which arises because the interaction involves the spin 1/2 components of the off shell ψ^μ . In the context of chiral perturbation theory (see Ref. [142] for a review), one can absorb z in the Wilson coefficients of certain contact terms via a field redefinition [143–145], rendering it redundant. In supergravity theories (see Ref. [146] for a review), the elementary spin 3/2 fermion known as the gravitino obtains its mass from a spin 1/2 fermion known as the Goldstino via the super-Higgs mechanism [147], so that it has a spin 1/2 component even when on shell, and one finds for its couplings to fermions and pseudoscalars that $A/2(1 + 4z) + z = -1/2$ [148,149]. In this work, we will not concern ourselves with the details of the off shell parameters because we are only interested in on shell composite spin 3/2 fermions; for the remainder of this work, we set $\theta_{\mu\nu} = g_{\mu\nu}$ similar to Ref. [134].

Since \mathcal{B} fields carries a Lorentz index and due to the constraint in Eq. (A2), the higher-dimensional operators destabilizing the dark baryon must involve derivatives. This is why the lowest allowed operator dimension for DM decay starts at $d = 8$ compared to DM with spin 1/2, where, e.g., dimension 7 operators are possible [55]. Schematically, the leading operators are at dark quark level,

$$\frac{c_8^{(1)}}{\Lambda_{\text{UV}}^4} (qqq)_\mu \bar{L} (D^\mu \tilde{H}) + \text{H.c.}, \quad (\text{A10})$$

$$\frac{c_8^{(2)}}{\Lambda_{\text{UV}}^4} (qqq)_\mu [\gamma^\alpha, \gamma^\beta] \gamma^\mu N_{L,R} B_{\alpha\beta} + \text{H.c.}, \quad (\text{A11})$$

$$\frac{c_{10}}{\Lambda_{\text{UV}}^6} (qqq)_\mu N_{L,R} (\partial^\mu \bar{q} q) + \text{H.c.}, \quad (\text{A12})$$

$$\frac{c_{11}}{\Lambda_{\text{UV}}^7} (qqq)_\mu \bar{L} (\partial^\mu \bar{q} q) \tilde{H} + \text{H.c.}, \quad (\text{A13})$$

where D^μ denotes the gauge covariant derivative, $B^{\mu\nu}$ is the hypercharge field strength, and $(qqq)_\mu$ indicates that we need a spinor vector from the symmetric spin contraction of three dark quarks for spin 3/2. The operators in Eqs. (A10) and (A11) were already mentioned in Ref. [150] for the gravitino and in Ref. [134] for a general elementary spin 3/2 fermion.

- [1] A. Giarnetti, J. Herrero-Garcia, S. Marciano, D. Meloni, and D. Vatsyayan, Neutrino masses from new Weinberg-like operators: Phenomenology of TeV scalar multiplets, *J. High Energy Phys.* **05** (2024) 055.
- [2] R. N. Mohapatra and J. W. F. Valle, Neutrino mass and baryon number nonconservation in superstring models, *Phys. Rev. D* **34**, 1642 (1986).
- [3] M. C. Gonzalez-Garcia and J. W. F. Valle, Fast decaying neutrinos and observable flavor violation in a new class of Majoron models, *Phys. Lett. B* **216**, 360 (1989).
- [4] Y. Cai, J. Herrero-García, M. A. Schmidt, A. Vicente, and R. R. Volkas, From the trees to the forest: A review of radiative neutrino mass models, *Front. Phys.* **5**, 63 (2017).
- [5] P. Minkowski, $\mu \rightarrow e\gamma$ at a rate of one out of 10^9 muon decays?, *Phys. Lett.* **67B**, 421 (1977).
- [6] T. Yanagida, Horizontal gauge symmetry and masses of neutrinos, eConf **C7902131**, 95 (1979).
- [7] M. Gell-Mann, P. Ramond, and R. Slansky, Complex spinors and unified theories, eConf **C790927**, 315 (1979).
- [8] S. L. Glashow, The future of elementary particle physics, *NATO Sci. Ser. B* **61**, 687 (1980).
- [9] T. Yanagida, Horizontal symmetry and masses of neutrinos, *Prog. Theor. Phys.* **64**, 1103 (1980).
- [10] R. N. Mohapatra and G. Senjanović, Neutrino mass and spontaneous parity nonconservation, *Phys. Rev. Lett.* **44**, 912 (1980).
- [11] R. Foot, H. Lew, X. G. He, and G. C. Joshi, Seesaw neutrino masses induced by a triplet of leptons, *Z. Phys. C* **44**, 441 (1989).
- [12] E. Witten, New issues in manifolds of SU(3) holonomy, *Nucl. Phys.* **B268**, 79 (1986).
- [13] E. Ma, Radiative inverse seesaw mechanism for nonzero neutrino mass, *Phys. Rev. D* **80**, 013013 (2009).
- [14] A. Ahriche, S. M. Boucenna, and S. Nasri, Dark radiative inverse seesaw mechanism, *Phys. Rev. D* **93**, 075036 (2016).
- [15] L. Coito, C. Faubel, J. Herrero-García, A. Santamaria, and A. Titov, Sterile neutrino portals to Majorana dark matter: Effective operators and UV completions, *J. High Energy Phys.* **08** (2022) 085.
- [16] S. D. Thomas and R.-M. Xu, Light neutrinos from the quark condensate, *Phys. Lett. B* **284**, 341 (1992).
- [17] L. E. Ibanez, F. Marchesano, and R. Rabadan, Getting just the standard model at intersecting branes, *J. High Energy Phys.* **11** (2001) 002.
- [18] J. McDonald, Nucleosynthesis bounds on small Dirac neutrino masses due to chiral symmetry breaking, [arXiv: hep-ph/9610324](https://arxiv.org/abs/hep-ph/9610324).
- [19] H. Davoudiasl and L. L. Everett, Implications of neutrino mass generation from QCD confinement, *Phys. Lett. B* **634**, 55 (2006).
- [20] A. Babič, S. Kovalenko, M. I. Krivoruchenko, and F. Šimkovic, Quark condensate seesaw mechanism for neutrino mass, *Phys. Rev. D* **103**, 015007 (2021).
- [21] H. Davoudiasl, I. M. Lewis, and M. Sullivan, Exploring strange origin of Dirac neutrino masses at hadron colliders, *Phys. Rev. D* **105**, 075017 (2022).
- [22] A. Bazavov *et al.* (Fermilab Lattice, MILC, and TUMQCD Collaborations), Up-, down-, strange-, charm-, and bottom-quark masses from four-flavor lattice QCD, *Phys. Rev. D* **98**, 054517 (2018).
- [23] C. Alexandrou, J. Finkenrath, L. Funcke, K. Jansen, B. Kostrzewa, F. Pittler, and C. Urbach, Ruling out the massless up-quark solution to the strong CP problem by computing the topological mass contribution with lattice QCD, *Phys. Rev. Lett.* **125**, 232001 (2020).
- [24] T. Hur, D.-W. Jung, P. Ko, and J. Y. Lee, Electroweak symmetry breaking and cold dark matter from strongly interacting hidden sector, *Phys. Lett. B* **696**, 262 (2011).
- [25] J. Kubo, K. S. Lim, and M. Lindner, Electroweak symmetry breaking via QCD, *Phys. Rev. Lett.* **113**, 091604 (2014).
- [26] T. Hur and P. Ko, Scale invariant extension of the Standard Model with strongly interacting hidden sector, *Phys. Rev. Lett.* **106**, 141802 (2011).
- [27] J. Gluza, On teraelectronvolt Majorana neutrinos, *Acta Phys. Pol. B* **33**, 1735 (2002).
- [28] J. Kersten and A. Y. Smirnov, Right-handed neutrinos at CERN LHC and the mechanism of neutrino mass generation, *Phys. Rev. D* **76**, 073005 (2007).
- [29] M. Aoki, V. Brdar, and J. Kubo, Heavy dark matter, neutrino masses, and Higgs naturalness from a strongly interacting hidden sector, *Phys. Rev. D* **102**, 035026 (2020).
- [30] M. Aoki, J. Kubo, and J. Yang, Inflation and dark matter after spontaneous Planck scale generation by hidden chiral symmetry breaking, *J. Cosmol. Astropart. Phys.* **01** (2022) 005.
- [31] I. Brivio and M. Trott, Radiatively generating the Higgs potential and electroweak scale via the seesaw mechanism, *Phys. Rev. Lett.* **119**, 141801 (2017).
- [32] I. Brivio and M. Trott, Examining the neutrino option, *J. High Energy Phys.* **02** (2019) 107.
- [33] E. Hall, T. Konstandin, R. McGehee, H. Murayama, and G. Servant, Baryogenesis from a dark first-order phase transition, *J. High Energy Phys.* **04** (2020) 042.
- [34] E. Hall, T. Konstandin, R. McGehee, and H. Murayama, Asymmetric matter from a dark first-order phase transition, *Phys. Rev. D* **107**, 055011 (2023).
- [35] E. Hall, R. McGehee, H. Murayama, and B. Suter, Asymmetric dark matter may not be light, *Phys. Rev. D* **106**, 075008 (2022).
- [36] G. B. Gelmini and M. Roncadelli, Left-handed neutrino mass scale and spontaneously broken lepton number, *Phys. Lett.* **99B**, 411 (1981).
- [37] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Are there real Goldstone bosons associated with broken lepton number?, *Phys. Lett.* **98B**, 265 (1981).
- [38] A. Mitridate, M. Redi, J. Smirnov, and A. Strumia, Dark matter as a weakly coupled dark baryon, *J. High Energy Phys.* **10** (2017) 210.
- [39] G. Lazarides, Q. Shafi, and C. Wetterich, Proton lifetime and fermion masses in an SO(10) model, *Nucl. Phys.* **B181**, 287 (1981).
- [40] J. Schechter and J. W. F. Valle, Neutrino masses in SU(2) \times U(1) theories, *Phys. Rev. D* **22**, 2227 (1980).
- [41] R. N. Mohapatra and G. Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, *Phys. Rev. D* **23**, 165 (1981).
- [42] T. P. Cheng and L.-F. Li, Neutrino masses, mixings, and oscillations in SU(2) \times U(1) models of electroweak interactions, *Phys. Rev. D* **22**, 2860 (1980).

- [43] C. Wetterich, Neutrino masses and the scale of B-L violation, *Nucl. Phys.* **B187**, 343 (1981).
- [44] E. Ma, Naturally small seesaw neutrino mass with no new physics beyond the TeV scale, *Phys. Rev. Lett.* **86**, 2502 (2001).
- [45] S. M. Davidson and H. E. Logan, Dirac neutrinos from a second Higgs doublet, *Phys. Rev. D* **80**, 095008 (2009).
- [46] S. Centelles Chuliá, R. Srivastava, and A. Vicente, The inverse seesaw family: Dirac and Majorana, *J. High Energy Phys.* **03** (2021) 248.
- [47] P. Sikivie, Of axions, domain walls and the early universe, *Phys. Rev. Lett.* **48**, 1156 (1982).
- [48] J. R. Espinosa, T. Konstandin, and F. Riva, Strong electroweak phase transitions in the standard model with a singlet, *Nucl. Phys.* **B854**, 592 (2012).
- [49] J. Preskill, S. P. Trivedi, F. Wilczek, and M. B. Wise, Cosmology and broken discrete symmetry, *Nucl. Phys.* **B363**, 207 (1991).
- [50] A. Francis, R. J. Hudspith, R. Lewis, and S. Tulin, Dark matter from strong dynamics: The minimal theory of dark baryons, *J. High Energy Phys.* **12** (2018) 118.
- [51] M. Della Morte, B. Jäger, F. Sannino, J. T. Tsang, and F. P. G. Ziegler, Spectrum of QCD with one flavor: A window for supersymmetric dynamics, *Phys. Rev. D* **107**, 114506 (2023).
- [52] M. Gell-Mann, R. J. Oakes, and B. Renner, Behavior of current divergences under $SU(3) \times SU(3)$, *Phys. Rev.* **175**, 2195 (1968).
- [53] T. Hambye, M. Hufnagel, and M. Lucca, Cosmological constraints on the decay of heavy relics into neutrinos, *J. Cosmol. Astropart. Phys.* **05** (2022) 033.
- [54] O. Antipin, M. Redi, A. Strumia, and E. Vigiani, Accidental composite dark matter, *J. High Energy Phys.* **07** (2015) 039.
- [55] R. Garani, M. Redi, and A. Tesi, Dark QCD matters, *J. High Energy Phys.* **12** (2021) 139.
- [56] W. Ochs, The status of glueballs, *J. Phys. G* **40**, 043001 (2013).
- [57] T. Cohen, F. J. Llanes-Estrada, J. R. Pelaez, and J. Ruiz de Elvira, Nonordinary light meson couplings and the $1/N_c$ expansion, *Phys. Rev. D* **90**, 036003 (2014).
- [58] D. Buttazzo, L. Di Luzio, G. Landini, A. Strumia, and D. Teresi, Dark matter from self-dual gauge/Higgs dynamics, *J. High Energy Phys.* **10** (2019) 067.
- [59] G. Landini and J.-W. Wang, Dark matter in scalar $Sp(\mathcal{N})$ gauge dynamics, *J. High Energy Phys.* **06** (2020) 167.
- [60] C. Gross, S. Karamitsos, G. Landini, and A. Strumia, Gravitational vector dark matter, *J. High Energy Phys.* **03** (2021) 174.
- [61] R. Contino, A. Mitridate, A. Podo, and M. Redi, Gluequark dark matter, *J. High Energy Phys.* **02** (2019) 187.
- [62] G. F. Giudice, E. W. Kolb, and A. Riotto, Largest temperature of the radiation era and its cosmological implications, *Phys. Rev. D* **64**, 023508 (2001).
- [63] E. W. Kolb, A. Notari, and A. Riotto, On the reheating stage after inflation, *Phys. Rev. D* **68**, 123505 (2003).
- [64] C. Gross, A. Mitridate, M. Redi, J. Smirnov, and A. Strumia, Cosmological abundance of colored relics, *Phys. Rev. D* **99**, 016024 (2019).
- [65] G. Steigman, B. Dasgupta, and J. F. Beacom, Precise relic WIMP abundance and its impact on searches for dark matter annihilation, *Phys. Rev. D* **86**, 023506 (2012).
- [66] J. M. Cline, Z. Liu, G. D. Moore, and W. Xue, Composite strongly interacting dark matter, *Phys. Rev. D* **90**, 015023 (2014).
- [67] A. Borriello and P. Salucci, The dark matter distribution in disk galaxies, *Mon. Not. R. Astron. Soc.* **323**, 285 (2001).
- [68] F. Donato, G. Gentile, P. Salucci, C. Frigerio Martins, M. I. Wilkinson, G. Gilmore, E. K. Grebel, A. Koch, and R. Wyse, A constant dark matter halo surface density in galaxies, *Mon. Not. R. Astron. Soc.* **397**, 1169 (2009).
- [69] W. J. G. de Blok, The core-cusp problem, *Adv. Astron.* **2010**, 789293 (2010).
- [70] W. J. G. de Blok and A. Bosma, High-resolution rotation curves of low surface brightness galaxies, *Astron. Astrophys.* **385**, 816 (2002).
- [71] M. Boylan-Kolchin, J. S. Bullock, and M. Kaplinghat, Too big to fail? The puzzling darkness of massive Milky Way subhaloes, *Mon. Not. R. Astron. Soc.* **415**, L40 (2011).
- [72] M. Boylan-Kolchin, J. S. Bullock, and M. Kaplinghat, The Milky Way's bright satellites as an apparent failure of Λ CDM, *Mon. Not. R. Astron. Soc.* **422**, 1203 (2012).
- [73] J. Miralda-Escude, A test of the collisional dark matter hypothesis from cluster lensing, *Astrophys. J.* **564**, 60 (2002).
- [74] S. W. Randall, M. Markevitch, D. Clowe, A. H. Gonzalez, and M. Bradac, Constraints on the self-interaction cross-section of dark matter from numerical simulations of the merging galaxy cluster 1E 0657-56, *Astrophys. J.* **679**, 1173 (2008).
- [75] D. Pappadopulo, J. T. Ruderman, and G. Trevisan, Dark matter freeze-out in a nonrelativistic sector, *Phys. Rev. D* **94**, 035005 (2016).
- [76] M. Farina, D. Pappadopulo, J. T. Ruderman, and G. Trevisan, Phases of cannibal dark matter, *J. High Energy Phys.* **12** (2016) 039.
- [77] L. Morrison, S. Profumo, and D. J. Robinson, Large N -nightmare dark matter, *J. Cosmol. Astropart. Phys.* **05** (2021) 058.
- [78] S. Davidson, E. Nardi, and Y. Nir, Leptogenesis, *Phys. Rep.* **466**, 105 (2008).
- [79] G. F. Giudice, M. Peloso, A. Riotto, and I. Tkachev, Production of massive fermions at preheating and leptogenesis, *J. High Energy Phys.* **08** (1999) 014.
- [80] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); **652**, C4(E) (2021).
- [81] K. Agashe, P. Du, M. Ekhterachian, C. S. Fong, S. Hong, and L. Vecchi, Natural seesaw and leptogenesis from hybrid of high-scale type I and TeV-scale inverse, *J. High Energy Phys.* **04** (2019) 029.
- [82] A. Pilaftsis and T. E. J. Underwood, Resonant leptogenesis, *Nucl. Phys.* **B692**, 303 (2004).
- [83] A. Pilaftsis, Resonant tau-leptogenesis with observable lepton number violation, *Phys. Rev. Lett.* **95**, 081602 (2005).
- [84] A. Pilaftsis and T. E. J. Underwood, Electroweak-scale resonant leptogenesis, *Phys. Rev. D* **72**, 113001 (2005).

- [85] P.D. Bolton, F.F. Deppisch, and P.S. Bhupal Dev, Neutrinoless double beta decay versus other probes of heavy sterile neutrinos, *J. High Energy Phys.* **03** (2020) 170.
- [86] A. Abada, C. Biggio, F. Bonnet, M.B. Gavela, and T. Hambye, Low energy effects of neutrino masses, *J. High Energy Phys.* **12** (2007) 061.
- [87] E. Fernandez-Martinez, J. Hernandez-Garcia, and J. Lopez-Pavon, Global constraints on heavy neutrino mixing, *J. High Energy Phys.* **08** (2016) 033.
- [88] M. Blennow, P. Coloma, E. Fernandez-Martinez, J. Hernandez-Garcia, and J. Lopez-Pavon, Non-unitarity, sterile neutrinos, and non-standard neutrino interactions, *J. High Energy Phys.* **04** (2017) 153.
- [89] S. Antusch, E. Cazzato, and O. Fischer, Sterile neutrino searches at future e^-e^+ , pp , and e^-p colliders, *Int. J. Mod. Phys. A* **32**, 1750078 (2017).
- [90] S. Antusch, E. Cazzato, and O. Fischer, Displaced vertex searches for sterile neutrinos at future lepton colliders, *J. High Energy Phys.* **12** (2016) 007.
- [91] A. Blondel *et al.*, Searches for long-lived particles at the future FCC-ee, *Front. Phys.* **10**, 967881 (2022).
- [92] A.M. Abdullahi *et al.*, The present and future status of heavy neutral leptons, *J. Phys. G* **50**, 020501 (2023).
- [93] A.M. Baldini *et al.* (MEG Collaboration), Search for the lepton flavour violating decay $\mu^+ \rightarrow e^+\gamma$ with the full dataset of the MEG experiment, *Eur. Phys. J. C* **76**, 434 (2016).
- [94] A.M. Baldini *et al.* (MEG II Collaboration), The design of the MEG II experiment, *Eur. Phys. J. C* **78**, 380 (2018).
- [95] K. Afanaciev *et al.* (MEG II Collaboration), A search for $\mu^+ \rightarrow e^+\gamma$ with the first dataset of the MEG II experiment, *Eur. Phys. J. C* **84**, 216 (2024).
- [96] C. Dohmen *et al.* (SINDRUM II Collaboration), Test of lepton flavor conservation in $\mu \rightarrow e$ conversion on titanium, *Phys. Lett. B* **317**, 631 (1993).
- [97] A. de Gouvea and S. Gopalakrishna, Low-energy neutrino Majorana phases and charged-lepton electric dipole moments, *Phys. Rev. D* **72**, 093008 (2005).
- [98] A. Abada and T. Toma, Electric dipole moments of charged leptons with sterile fermions, *J. High Energy Phys.* **02** (2016) 174.
- [99] ACME Collaboration, Improved limit on the electric dipole moment of the electron, *Nature (London)* **562**, 355 (2018).
- [100] A. Abada and T. Toma, Electron electric dipole moment in inverse seesaw models, *J. High Energy Phys.* **08** (2016) 079.
- [101] G.W. Bennett *et al.* (Muon $g-2$ Collaboration), An improved limit on the muon electric dipole moment, *Phys. Rev. D* **80**, 052008 (2009).
- [102] Y. Ema, T. Gao, and M. Pospelov, Improved indirect limits on muon electric dipole moment, *Phys. Rev. Lett.* **128**, 131803 (2022).
- [103] A. Adelman *et al.*, Search for a muon EDM using the frozen-spin technique, [arXiv:2102.08838](https://arxiv.org/abs/2102.08838).
- [104] M. Sakurai *et al.*, muEDM: Towards a search for the muon electric dipole moment at PSI using the frozen-spin technique, *J. Phys. Soc. Jpn. Conf. Proc.* **37**, 020604 (2022).
- [105] K.S. Khaw *et al.* (Muon EDM Initiative Collaboration), Search for the muon electric dipole moment using frozen-spin technique at PSI, *Proc. Sci. NuFact2021* (2022) 136 [[arXiv:2201.08729](https://arxiv.org/abs/2201.08729)].
- [106] J.L. Feng, K.T. Matchev, and Y. Shadmi, Muon dipole moment experiments: Interpretation and prospects, *eConf C010630*, P307 (2001).
- [107] J. a. P. Pinheiro, C. A. de S. Pires, F. S. Queiroz, and Y. S. Villamizar, Confronting the inverse seesaw mechanism with the recent muon $g-2$ result, *Phys. Lett. B* **823**, 136764 (2021).
- [108] B. Abi *et al.* (Muon $g-2$ Collaboration), Measurement of the positive muon anomalous magnetic moment to 0.46 ppm, *Phys. Rev. Lett.* **126**, 141801 (2021).
- [109] P.B. Denton and J. Gehrlein, A survey of neutrino flavor models and the neutrinoless double beta decay funnel, *Phys. Rev. D* **109**, 055028 (2024).
- [110] Y. Bai, A. J. Long, and S. Lu, Dark quark nuggets, *Phys. Rev. D* **99**, 055047 (2019).
- [111] A. Das, S. Jana, S. Mandal, and S. Nandi, Probing right handed neutrinos at the LHeC and lepton colliders using fat jet signatures, *Phys. Rev. D* **99**, 055030 (2019).
- [112] V. Berezhinsky, M. Kachelriess, and S. Ostapchenko, Electroweak jet cascading in the decay of superheavy particles, *Phys. Rev. Lett.* **89**, 171802 (2002).
- [113] M. Kachelriess, P.D. Serpico, and M.A. Solberg, On the role of electroweak bremsstrahlung for indirect dark matter signatures, *Phys. Rev. D* **80**, 123533 (2009).
- [114] P. Ciafaloni, D. Comelli, A. Riotto, F. Sala, A. Strumia, and A. Urbano, Weak corrections are relevant for dark matter indirect detection, *J. Cosmol. Astropart. Phys.* **03** (2011) 019.
- [115] C. El Aisati, M. Gustafsson, and T. Hambye, New search for monochromatic neutrinos from dark matter decay, *Phys. Rev. D* **92**, 123515 (2015).
- [116] M.G. Aartsen *et al.* (IceCube Collaboration), Energy reconstruction methods in the IceCube neutrino telescope, *J. Instrum.* **9**, P03009 (2014).
- [117] C. A. Argüelles, A. Diaz, A. Kheirandish, A. Olivares-Del-Campo, I. Safa, and A. C. Vincent, Dark matter annihilation to neutrinos, *Rev. Mod. Phys.* **93**, 035007 (2021).
- [118] ANTARES Collaboration, Search of dark matter annihilation in the Galactic centre using the ANTARES neutrino telescope, *J. Cosmol. Astropart. Phys.* **10** (2015) 068.
- [119] M.G. Aartsen *et al.* (IceCube Collaboration), Neutrino astronomy with the next generation IceCube neutrino observatory, [arXiv:1911.02561](https://arxiv.org/abs/1911.02561).
- [120] M.G. Aartsen *et al.* (IceCube Collaboration), IceCube-Gen2: A vision for the future of neutrino astronomy in antarctica, [arXiv:1412.5106](https://arxiv.org/abs/1412.5106).
- [121] M. Agostini *et al.* (P-ONE Collaboration), The Pacific Ocean neutrino experiment, *Nat. Astron.* **4**, 913 (2020).
- [122] S. Adrian-Martinez *et al.* (KM3Net Collaboration), Letter of intent for KM3NeT 2.0, *J. Phys. G* **43**, 084001 (2016).
- [123] S. Aiello *et al.* (KM3NeT Collaboration), Sensitivity of the KM3NeT/ARCA neutrino telescope to point-like neutrino sources, *Astropart. Phys.* **111**, 100 (2019).
- [124] M. Ackermann *et al.* (Fermi-LAT Collaboration), Searching for dark matter annihilation from Milky Way dwarf

- spheroidal galaxies with six years of Fermi Large Area Telescope data, *Phys. Rev. Lett.* **115**, 231301 (2015).
- [125] H. Abdallah *et al.* (H.E.S.S. Collaboration), Search for dark matter annihilations towards the inner Galactic halo from 10 years of observations with H.E.S.S., *Phys. Rev. Lett.* **117**, 111301 (2016).
- [126] F. S. Queiroz, C. E. Yaguna, and C. Weniger, Gamma-ray limits on neutrino lines, *J. Cosmol. Astropart. Phys.* **05** (2016) 050.
- [127] B. Batell, T. Han, and B. Shams Es Haghi, Indirect detection of neutrino portal dark matter, *Phys. Rev. D* **97**, 095020 (2018).
- [128] A. Morselli (CTA Consortium Collaboration), The dark matter programme of the Cherenkov Telescope Array, *Proc. Sci. ICRC2017* (**2018**) 921 [arXiv:1709.01483].
- [129] S. Coleman, Why there is nothing rather than something: A theory of the cosmological constant, *Nucl. Phys.* **B310**, 643 (1988).
- [130] S. B. Giddings and A. Strominger, Loss of incoherence and determination of coupling constants in quantum gravity, *Nucl. Phys.* **B307**, 854 (1988).
- [131] G. Gilbert, Wormhole-induced proton decay, *Nucl. Phys.* **B328**, 159 (1989).
- [132] R. Kallosh, A. D. Linde, D. A. Linde, and L. Susskind, Gravity and global symmetries, *Phys. Rev. D* **52**, 912 (1995).
- [133] Y. Mambrini, S. Profumo, and F. S. Queiroz, Dark matter and global symmetries, *Phys. Lett. B* **760**, 807 (2016).
- [134] M. A. G. Garcia, Y. Mambrini, K. A. Olive, and S. Verner, Case for decaying spin-3/2 dark matter, *Phys. Rev. D* **102**, 083533 (2020).
- [135] M. G. Aartsen *et al.* (IceCube Collaboration), Atmospheric and astrophysical neutrinos above 1 TeV interacting in IceCube, *Phys. Rev. D* **91**, 022001 (2015).
- [136] W. Rarita and J. Schwinger, On a theory of particles with half integral spin, *Phys. Rev.* **60**, 61 (1941).
- [137] S. Weinberg, *The Quantum Theory of Fields. Vol. 1: Foundations* (Cambridge University Press, Cambridge, England, 2005), 10.1017/CBO9781139644167.
- [138] H. Haberzettl, Propagation of a massive spin 3/2 particle, arXiv:nucl-th/9812043.
- [139] L. M. Nath, B. Etemadi, and J. D. Kimel, Uniqueness of the interaction involving spin 3/2 particles, *Phys. Rev. D* **3**, 2153 (1971).
- [140] M. Benmerrouche, R. M. Davidson, and N. C. Mukhopadhyay, Problems of describing spin 3/2 baryon resonances in the effective Lagrangian theory, *Phys. Rev. C* **39**, 2339 (1989).
- [141] S. Kamefuchi, L. O’Raifeartaigh, and A. Salam, Change of variables and equivalence theorems in quantum field theories, *Nucl. Phys.* **28**, 529 (1961).
- [142] S. Scherer, Introduction to chiral perturbation theory, *Adv. Nucl. Phys.* **27**, 277 (2003).
- [143] H.-B. Tang and P. J. Ellis, Redundance of delta isobar parameters in effective field theories, *Phys. Lett. B* **387**, 9 (1996).
- [144] P. J. Ellis and H.-B. Tang, Pion—nucleon scattering at low-energies, *Phys. Rev. C* **56**, 3363 (1997).
- [145] H. Krebs, E. Epelbaum, and U. G. Meissner, Redundancy of the off-shell parameters in chiral effective field theory with explicit spin-3/2 degrees of freedom, *Phys. Lett. B* **683**, 222 (2010).
- [146] H. P. Nilles, Supersymmetry, supergravity and particle physics, *Phys. Rep.* **110**, 1 (1984).
- [147] E. Cremmer, B. Julia, J. Scherk, P. van Nieuwenhuizen, S. Ferrara, and L. Girardello, Super-Higgs effect in supergravity with general scalar interactions, *Phys. Lett.* **79B**, 231 (1978).
- [148] S. Ferrara, F. Gliozzi, J. Scherk, and P. Van Nieuwenhuizen, Matter couplings in supergravity theory, *Nucl. Phys.* **B117**, 333 (1976).
- [149] A. Das, M. Fischler, and M. Roček, Super-Higgs effect in a new class of scalar models and a model of super QED, *Phys. Rev. D* **16**, 3427 (1977).
- [150] E. Dudas, T. Gherghetta, K. Kaneta, Y. Mambrini, and K. A. Olive, Gravitino decay in high scale supersymmetry with R-parity violation, *Phys. Rev. D* **98**, 015030 (2018).