Bound-state dark matter and Dirac neutrino masses

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We propose a simple theory for the idea that cosmological dark matter (DM) may be present today mainly in the form of stable neutral hadronic thermal relics. In our model, neutrino masses arise radiatively from the exchange of colored DM constituents, giving a common origin for both dark matter and neutrino mass. The exact conservation of $B - L$ symmetry ensures dark matter stability and the Dirac nature of neutrinos. The theory can be falsified by dark matter nuclear recoil direct detection experiments, leading also to possible signals at a next generation hadron collider.

I. INTRODUCTION

The common lore concerning particle dark matter candidates has long been that they must be electrically neutral and carry no color. This view has been challenged by the authors in Ref. [1], who suggested that dark matter (DM) may be the lightest hadron made of two stable color octet Dirac fermions $\mathcal{Q}$ with mass below 10 TeV. The interest in this idea may go well beyond QCD, since analogous bound-state DM candidates emerge in models with a new confining hypercolor interaction [2].

We argue that neutrino mass and cosmological dark matter may have a common origin, with the underlying DM physics acting as messenger of neutrino mass generation [3–6]. We propose a simple implementation of the bound-state dark matter scenario in which DM constituents induce calculable neutrino mass at the radiative level. Dark matter is a QCD bound state $\mathcal{Q}\mathcal{Q}$ stabilized by the same conserved $B - L$ symmetry associated to the Dirac nature of neutrinos [7].

In addition to the heavy Dirac fermion, our model introduces extra scalars, in order to ensure that at least two neutrino masses are nonzero, required by the neutrino oscillation data in order to account both for solar and atmospheric mass scales [8,9]. This simple picture can account for current neutrino oscillation and dark matter

phenomena, and can be falsified relatively soon, in nuclear recoil studies at XENON1T [10]. Moreover, the extra colored states, including scalar bosons, may lead to new phenomena at a next generation hadron collider.

II. THE MODEL

As a theory preliminary, we recall that, within the type-I seesaw mechanism with a single right-handed neutrino, two neutrinos remain massless after the seesaw [11]. This degeneracy is lifted by calculable loop corrections.[1] Here we propose a variant radiative seesaw scheme, in which a single colored fermion Dirac messenger $\mathcal{Q}$ is introduced, charged under the $U(1)_{B-L}$ symmetry, plus two sets of colored scalars, labeled by $a = 1, 2$; see Table I.

Apart from the right-handed neutrinos, all of the new particles are colored. For definiteness, we assign them to the octet $SU(3)_c$ representation.[2] We also impose that the $B - L$ symmetry holds, together with a Z_2 symmetry. The former ensures that $\mathcal{Q}$ has only a Dirac-type mass term, while the latter forbids the tree-level Dirac neutrino mass terms from $\overline{\nu_{Ri}}\tilde{H}^\dagger L_j$. The Lagragian contains the following new terms (summation is implied over repeated indices, and trace over $\mathbf{N}_c = \mathbf{8}$ is implicit.)

$$\mathcal{L} \supset -[h_i^a \overline{\mathcal{Q}_R}\tilde{\eta}_a^\dagger L_i + M_\mathcal{Q} \overline{\mathcal{Q}_R}\mathcal{Q}_L + y_i^a \overline{\nu_{Ri}}\sigma_a^* \mathcal{Q}_L + \text{H.c.}]$$
$$- \mathcal{V}(H, \eta_a, \sigma_a). \tag{1}$$

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[1]An analogous situation happens in supersymmetry with bilinear breaking of R parity [12–14], which induces the solar scale at one-loop, once atmospheric scale is taken as tree level input.

[2]Notice, however, that our neutrino mass discussion also holds if they had different $SU(3)_c$ transformation properties.

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TABLE I. Left-handed fermions and scalars. When the vector-like quark $\mathcal{Q}$ is taken as a color octet, $N_c = 8$, one can realize bound-state dark matter scenario proposed in Ref. [1].

Particles	$U(1)_{B-L}$	$(SU(3)_c, SU(2)_L)_Y$	Z_2
$Q_i = (u_L d_L)_i^T$	$+1/3$	$(\mathbf{3}, \mathbf{2})_{1/6}$	$+$
$\overline{u_{Ri}}$	$-1/3$	$(\bar{\mathbf{3}}, \mathbf{1})_{-2/3}$	$+$
$\overline{d_{Ri}}$	$-1/3$	$(\bar{\mathbf{3}}, \mathbf{1})_{1/3}$	$+$
$L_i = (\nu_L e_L)_i^T$	-1	$(\mathbf{1}, \mathbf{2})_{-1/2}$	$+$
$\overline{e_{Ri}}$	$+1$	$(\mathbf{1}, \mathbf{1})_1$	$+$
$\overline{\nu_{Ri}}$	$+1$	$(\mathbf{1}, \mathbf{1})_0$	$-$
$\mathcal{Q}_L$	$-r$	$(N_c, \mathbf{1})_0$	$+$
$\overline{\mathcal{Q}_R}$	r	$(N_c, \mathbf{1})_0$	$+$
H	0	$(\mathbf{1}, \mathbf{2})_{1/2}$	$+$
σ_a	$1-r$	$(N_c, \mathbf{1})_0$	$-$
η_a	$1-r$	$(N_c, \mathbf{2})_{1/2}$	$+$

The condition $r \neq 1$ forbids Higgs-like Yukawa couplings of η^α to the Standard Model fermions (for $r = 1$ one would need an additional Z_2 symmetry, as in [15]). The new part of the scalar potential can be cast as

$$\mathcal{V}(H, \eta_a, \sigma_a) = \mathcal{V}(\eta_a) + \mathcal{V}(\sigma_a) + \mathcal{V}(\eta_a, \sigma_a) + \mathcal{V}(H, \eta_a) + \mathcal{V}(H, \sigma_a), \quad (2)$$

where the various terms in the Higgs potential are

$$\mathcal{V}(\eta_a, \sigma_a) = \kappa^{ab} \text{Tr}(\sigma_a \eta_b^\dagger) H + \lambda_{\sigma\eta}^{abcd} \text{Tr}(\sigma_a^* \sigma_b) \text{Tr}(\eta_c^\dagger \eta_d)$$
$$+ \tilde{\lambda}_{\sigma\eta}^{abcd} \text{Tr}(\sigma_a^* \eta_b) \text{Tr}(\sigma_c \eta_d^\dagger) + \text{H.c.}, \quad (3)$$

$$\mathcal{V}(H, \eta_a) = \lambda_{3\eta H}^{ab}(H^\dagger H) \text{Tr}(\eta_a^\dagger \eta_b) + \lambda_{4\eta H}^{ab} \text{Tr}[(\eta_a^\dagger H)(H^\dagger \eta_b)]$$
$$+ \text{H.c.}, \quad (4)$$

$$\mathcal{V}(H, \sigma_a) = \lambda_{\sigma H}^{ab} \text{Tr}(\sigma_a^* \sigma_b) H^\dagger H + \text{H.c.}. \quad (5)$$

Since CP conservation is assumed, the CP-even and CP-odd scalars do not mix. Moreover, terms like $(\eta^\dagger H)^2$ are also forbidden and, as a consequence, the real and imaginary parts of the scalars with nonzero $B-L$ charges are degenerate. Note, however, that the cubic scalar coupling terms κ^{ab} breaking the Z_2 symmetry softly allow for the mixing between the σ_a and η_a [16]. At the end, the 4×4 mass matrices for the CP-odd and CP-even scalars are equal, since the κ^{ab} terms do not break such a degeneracy.

In order to illustrate the neutrino mass generation mechanism, we consider the following block-diagonal mass matrix for the CP-even scalars (in the basis $S_R = (\eta_{1R}^0, \sigma_{1R}^0, \eta_{2R}^0, \sigma_{2R}^0)^T$):
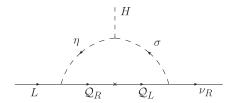

FIG. 1. $B - L$ flux in the one-loop Dirac neutrino mass.

$$\mathcal{M}_R^2 = \begin{pmatrix} (\mu_\eta^{11})^2 & \frac{\kappa_{11} v}{\sqrt{2}} & 0 & 0 \\ \frac{\kappa_{11} v}{\sqrt{2}} & (\mu_\sigma^{11})^2 & 0 & 0 \\ 0 & 0 & (\mu_\eta^{22})^2 & \frac{\kappa_{22} v}{\sqrt{2}} \\ 0 & 0 & \kappa_{22} \frac{v}{\sqrt{2}} & (\mu_\sigma^{22})^2 \end{pmatrix}. \quad (6)$$

Here we have used the parametrization $\eta_a = (\eta_a^+, (\eta_{aR}^0 + i\eta_{aI}^0)/\sqrt{2})^T$, $\sigma_a = (\sigma_{aR}^0 + i\sigma_{aI}^0)/\sqrt{2}$ and $H = (0, (h + v)/\sqrt{2})^T$, with h denoting the standard model Higgs boson, and $v = 246$ GeV. The parameters $\mu_{\eta(\sigma)}^{ab}$ are the quadratic mass terms after electroweak symmetry breaking in $\mathcal{V}(\eta_a)$ and $\mathcal{V}(\sigma_a)$.

Since the tree-level Dirac mass term is forbidden by symmetry, calculable neutrino masses are generated at one–loop order, by the Feynman diagram displayed in Fig. 1. One finds the following effective mass matrix

$$(\mathcal{M}_\nu)_{ij} = N_c \frac{M_\mathcal{Q}}{64\pi^2} \sum_{a=1}^2 h_i^a y_j^a \frac{\sqrt{2}\kappa_{aa} v}{m_{S_{2R}^a}^2 - m_{S_{1R}^a}^2}$$
$$\times \left[F\left(\frac{m_{S_{2R}^a}^2}{M_\mathcal{Q}^2}\right) - F\left(\frac{m_{S_{1R}^a}^2}{M_\mathcal{Q}^2}\right) \right] + (R \to I) \quad (7)$$

where $F(m_{S_\beta}^2/M_\mathcal{Q}^2) = m_{S_\beta}^2 \log(m_{S_\beta}^2/M_\mathcal{Q}^2)/(m_{S_\beta}^2 - M_\mathcal{Q}^2)$ and the $SU(3)_c$ color factor N_c is assumed to be 8, since the new particles running in the loop transform as octets. The four CP-even mass eigenstates are denoted as $S_{1R}^1, S_{2R}^1, S_{1R}^2, S_{2R}^2$, with a similar notation for the CP-odd ones. Notice that the above effective neutrino mass matrix in does not have the projective rank one structure $(\mathcal{M}_\nu)_{ij}^{\text{tree-level}} \propto h_i^a y_j^b$ characterizing the Dirac variant [7] of the "incomplete" tree-level seesaw mechanism, involving the exchange of a single heavy messenger, called (3,1) in Ref. [11].[3]

Here the effective one-loop induced neutrino mass matrix has in general rank two, implying two nonvanishing Dirac neutrino masses. As a simple numerical estimate, let's consider the case $\mu_\eta^{aa} = \mu_\sigma^{aa} \gg \sqrt{2}\kappa^{aa} v$. Taking $m_{S_{2R}^a}^2 - m_{S_{1R}^a}^2 = \sqrt{2}\kappa^{aa} v$ and $m_{S_{2R}^a}^2 + m_{S_{1R}^a}^2 = 2(\mu_\eta^{aa})^2$ one finds

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[3]Here one would recover the rank-one situation characterizing the type-I seesaw in the limit where one set of scalars decouples.

$$(\mathcal{M}_\nu)_{ij} = N_c \frac{M_\mathcal{Q}}{32\pi^2} \sqrt{2} \kappa^{aa} v \sum_{a=1}^{2} \frac{h_i^a y_j^a}{(\mu_\eta^{aa})^2 - M_\mathcal{Q}^2}$$

$$\times \left[1 - \frac{M_\mathcal{Q}^2}{(\mu_\eta^{aa})^2 - M_\mathcal{Q}^2} \log\left(\frac{(\mu_\eta^{aa})^2}{M_\mathcal{Q}^2}\right) \right],$$

so that if $(\mu_\eta^{aa})^2 \gg M_\mathcal{Q}^2$ one has

$$(\mathcal{M}_\nu)_{ij} = N_c \frac{M_\mathcal{Q}}{32\pi^2} \sqrt{2} v \sum_{a=1}^{2} \kappa^{aa} \frac{h_i^a y_j^a}{(\mu_\eta^{aa})^2}$$

$$\sim 0.03 \text{ eV} \left(\frac{M_\mathcal{Q}}{9.5 \text{ TeV}}\right) \left(\frac{\kappa^{aa}}{1 \text{ GeV}}\right)$$

$$\times \left(\frac{50 \text{ TeV}}{\mu_\eta^{aa}}\right)^2 \left(\frac{h_i^a y_j^a}{10^{-6}}\right). \tag{8}$$

One sees that, indeed, small neutrino masses arise naturally by taking reasonable values for the Yukawa couplings, small value for the soft breaking parameter κ^{ab}, as well as sufficiently large values for the scalar masses. Notice that the smallness of κ^{ab} is natural in t'Hooft's sense [17], as the theory attains a larger symmetry when $\kappa^{ab} \to 0$, i.e., the smallness of neutrino mass is symmetry protected. Indeed, the naturalness of scotogenic mass generation has been established explicitly in the literature; see, for example, Ref. [18].

In short, concerning neutrino mass generation, our model provides a colored variant of the one suggested in [15]. However, our proposal provides the simplest implementation of the radiative Dirac seesaw mechanism, where the dark matter stability is genuinely associated to the Dirac nature of neutrinos. In fact, in our model, the scalars are charged under B-L, so that the stabilizing symmetry invoked in [15] is not required. Moreover, although the neutrino mass generation is similar in both models, the details of the associated physics differ substantially.

III. BOUND-STATE DARK MATTER STABILITY FROM DIRAC NEUTRINOS

The Dirac nature of neutrinos may ensure dark matter stability, as suggested in [7]. Here we clone this idea with the proposal that dark matter may be present today mainly in the form of stable neutral hadronic thermal relics. For definiteness we assume DM is a neutral bound-state of colored constituents, such as $\mathcal{QQ}$, where $\mathcal{Q}$ is a vectorlike color octet isosinglet fermion. It was claimed that a necessary and suficient condition for dark matter stability in this case is the presence of a global $U(1)_D$ dark baryon number, under which the $\mathcal{Q}$ is charged [1]. In our present model construction, the role of such apparently *ad hoc* symmetry is played by the usual $B - L$ symmetry present in the Standard Model. In fact, in our model dark matter stability, and the Dirac nature of the exotic fermion $\mathcal{Q}$ and
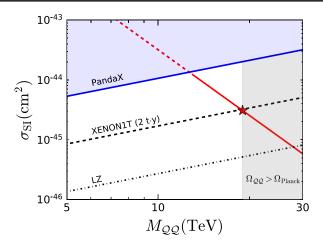

FIG. 2. The red line gives the spin independent cross section as a function of $M_{\mathcal{QQ}} = 2M_\mathcal{Q}$. The star denotes the mass required for a thermal bound state DM (19 TeV). Smaller values can be probed by direct searches, the blue line gives the current bound, while the black (dashed, dotted and dot-dashed) lines represent future sensitivities.

of the neutrinos are all equivalent, and result from $B - L$ conservation.

Hence, by construction, the new Yukawa interactions in Eq. (1) do not affect the stability of our bound state dark matter, since the colored scalars are heavier than $\mathcal{Q}$, and $\mathcal{QQ}$ annihilation processes mediated by η, σ, are all forbidden by the conserved $B - L$ symmetry.

An adequate thermal relic density of bound-state dark matter requires the lightest constituent vectorlike color octet Dirac fermion, $\mathcal{Q}$, to have a mass ≈ 9.5 TeV, so that the $\mathcal{QQ}$ hadron weighs approximately 19 TeV [1].

The set of scalars η^a, σ^a in Fig. 1 can be either neutral or charged under $B - L$, depending on the baryon number assignment. If $B - L$ neutral, both η^a and σ^a are expected to be unstable, decaying to quarks [19] and two gluons respectively [20,21]. In contrast, in the second case, if the lightest octet particle were a scalar, then it would be strictly stable.

Bound-state dark matter will impart nuclear recoil in underground dark matter search experiments. The spin-independent direct detection cross-section is given as

$$\sigma_{\text{SI}} \approx 2 \times 10^{-45} \text{ cm}^2 \left(\frac{20 \text{ TeV}}{M_{\mathcal{QQ}}}\right)^6 \frac{\Omega_{\mathcal{QQ}}}{\Omega_{\text{Planck}}}, \tag{9}$$

sharply correlated to the dark matter mass $M_{\mathcal{QQ}} = 2M_\mathcal{Q}$, as shown in the red line in Fig. 2.

In this figure, the star corresponds to the case where our bound-state DM makes up 100% of the cosmological dark matter. In the presence of an additional dark matter particle, such as the axion, bound-state dark matter masses below 19 TeV can be envisaged, as indicated by the red line. In this case, their contribution to the relic density will be correspondingly smaller, while the spin-independent cross

section would be correspondingly larger. The blue line represents the current limit of PandaX [22]. The black (dashed, and dot-dashed) lines represent the future sensitivities expected at XENON1T [10] and LZ [23]. On the other hand we note that, within the standard thermal cosmological scenario, DM masses above 19 TeV are ruled out by current observations by the Planck collaboration [24] (gray band).

Notice that the current LHC limit of 2 TeV (next section) implies that the cross section is always small enough so as to have the bound-state dark matter candidate reaching underground detectors. For more detailed discussion and general aspects of the cosmology of a stable colored relic, see [25].

IV. COLOR OCTETS AT HADRON COLLIDERS

In our model, the messengers of neutrino mass generation are the colored constituents of bound-state dark matter. Given enough energy, the $\mathcal{Q}$'s are copiously pair produced at hadron colliders, through the processes $q_i \bar{q}_i \to \mathcal{Q}\bar{\mathcal{Q}}$ and $gg \to \mathcal{Q}\bar{\mathcal{Q}}$, and are expected to hadronize. In contrast to WIMP dark matter scenarios, which engender only missing-energy signals, the bound-state dark matter scenario gives rise to very visible signals at hadron colliders, as they can form either neutral or charged bound states [1], e.g., neutral hybrid states $\mathcal{Q}g$ (detected as neutral hadrons, presumably stable) or charged $\mathcal{Q}q\bar{q}'$ states, or more exotic $\mathcal{Q}qqq$ states, expected to be long-lived on collider time scales.

Current LHC data place a limit to the fermion color octet mass, $M_{\mathcal{Q}} > 2$ TeV [26]. Since the cosmological relic abundance requires $M_{\mathcal{Q}} \approx 9.5$ TeV, this scenario offers an attractive benchmark for future collider experiments beyond the energies attainable at the LHC. In fact, from the estimate in Ref. [27] one finds that a hadron collider of at least 65 TeV center-of-mass energy would be required to probe the full cosmologically allowed range of masses of our bound-state DM model. This will allow a cross-check of the DM search results of XENON1T, expected quite soon, in just one year or so. Concerning the scalar messengers, we have two pairs of these, the $\sigma_a \sim (\mathbf{8}, 1, 0)$, which are singlets under $SU(2)_L$, and the η_a, which transform as weak doublets, $\eta_a \sim (\mathbf{8}, 2, 1/2)$. As color octets, these would also be copiously produced at a hadron collider of sufficient energy [28]. However, their masses are expected to lie well above the reach of the LHC. Moreover, in our model these scalars carry nontrivial $B - L$ charges; see Table I. This makes them relatively *inert* with respect to the standard model fermions. This, in addition to their heavy masses, makes them very difficult to probe directly.

V. LEPTON FLAVOR VIOLATION

Our model may also lead to indirect virtual effects, such as charged lepton flavor violation. For example, the Yukawa interactions in Eq. (1) lead to radiative lepton
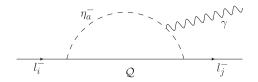

FIG. 3. Feynman diagram for the process $l_i \to l_j \gamma$.

flavor violation processes, as seen in Fig. 3, mediated by the charged scalar η_a^+.

The corresponding decay rate is given as [29]

$$\Gamma(l_i \to l_j \gamma) = \frac{e^2 m_i^5}{16\pi} \left[\sum_a N_c h_i^a h_j^{a*} \frac{i}{16\pi^2 M_{\eta_a^+}^2} \right.$$
$$\left. \times \left[\frac{-t^2 \log t}{2(t-1)^4} + \frac{2t^2 + 5t - 1}{12(t-1)^3} \right] \right]^2, \quad (10)$$

with $t = M_{\mathcal{Q}}^2 / M_{\eta^+}^2$ and $N_c = 8$. In the limit of heavy scalars, $t \to 0$, the decay width reads

$$\Gamma(l_i \to l_j \gamma) = \frac{e^2 m_i^5}{16\pi} \left[\sum_a N_c h_i^a h_j^{a*} \frac{i}{16\pi^2 M_{\eta_a^+}^2} \left[-\frac{1}{12} \right] \right]^2.$$

One sees that the current experimental constraint $\mathrm{BR}(\mu \to e\gamma) = \frac{\Gamma(\mu \to e\gamma)}{\Gamma_{\mathrm{total}}} < 5.7 \times 10^{-13}$, can be fulfilled provided

$$\left(\sum_a \frac{h_\mu^a h_e^{a*}}{M_{\eta_a^+}^2} \right)^2 \leq 5.7 \times 10^{-13} \frac{G_F^2}{\alpha_{\mathrm{EM}}} \frac{768\pi}{N_c^2}, \quad (11)$$

which leads to a relatively mild requirement,

$$\left| \sum_a h_\mu^a h_e^{a*} \left(\frac{50 \text{ TeV}}{M_{\eta_a^+}} \right)^2 \right| \lesssim 1.5. \quad (12)$$

VI. SUMMARY AND OUTLOOK

We have proposed a consistent viable theory for the recently proposed idea that the cosmological dark matter may be made up of stable colored relics forming neutral hadronic bound states of QCD. In our model, we have taken up at face value the suggestion in Ref. [1], employing an exotic vectorlike Dirac color octet fermion $\mathcal{Q}$ with mass below 10 TeV as the dark matter constituent. In our construction, dark matter and neutrino mass generation both have a common origin. Our minimum particle content leads to two nonzero neutrino masses, associated to the solar and atmospheric scales. Bound-state dark matter stability is directly associated with the Dirac nature of neutrinos, and reflects the presence of an underlying exact $B - L$ symmetry. The scheme can account for both neutrino physics and dark matter phenomena, within a consistent ultraviolet complete setup, free of Landau poles up to the

Planck scale, provided the scalars are heavy enough. Our model can be falsified relatively soon by dark matter searches, and could also be cross-checked later by a next generation hadron collider. Variants of our construction may be envisaged, in which the dark matter is bound by a new hypercolor interaction [30], instead of QCD, as suggested in Ref. [2].

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