



Bound-state dark matter with Majorana neutrinos

M. Reig^a, D. Restrepo^{b,c}, J.W.F. Valle^{a,*}, O. Zapata^b

^a AHEP Group, Institut de Física Corpuscular (CSIC-Universitat de València), Parc Científic de Paterna, C/Catedrático José Beltrán, 2 E-46980 Paterna (València), Spain

^b Instituto de Física, Universidad de Antioquia, Calle 70 No. 52-21, Apartado Aéreo 1226, Medellín, Colombia

^c Simons Associate at The Abdus Salam International Centre for Theoretical Physics (ICTP), Strada Costiera 11, 34151, Trieste, Italy

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ABSTRACT

We propose a simple scenario in which dark matter (DM) emerges as a stable neutral hadronic thermal relic, its stability following from an exact $U(1)_D$ symmetry. Neutrinos pick up radiatively induced Majorana masses from the exchange of colored DM constituents. There is a common origin for both dark matter and neutrino mass, with a lower bound for neutrinoless double beta decay. Direct DM searches at nuclear recoil experiments will test the proposal, which may also lead to other phenomenological signals at future hadron collider and lepton flavor violation experiments.

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1. Introduction

For a number of reasons, particle dark matter (DM) candidates are usually assumed to be electrically neutral and colorless. This has recently been challenged, however, in Ref. [1], where it was suggested that dark matter could be the lightest hadron made of two stable color octet Dirac fermions Q with mass around 12.5 TeV. Bound-state dark matter may also arise from new confining hypercolor interactions beyond QCD [2–4]. Such QCD-like bound state dark matter has been considered in many contexts [5], including asymmetric dark matter scenarios [6,7]. It can also emerge naturally in models of comprehensive unification [8]. A simple theory implementation of the bound-state dark matter scenario in which a QCD bound state $Q\bar{Q}$ is formed and its stability arises from the conservation of $B - L$ symmetry, has been proposed in [9]. In such picture the Diracness of neutrinos would be responsible for dark matter stability [10].

In this letter we propose an alternative bound-state dark matter construction in which $B - L$ symmetry is violated and neutrinos are Majorana fermions. In such scenario dark matter is stabilized by an extra continuous Abelian dark symmetry, $U(1)_D$, under which all standard model particles are assumed to be neutral. This leads to an economical alternative version of the bound-state DM picture in which neutrinos have radiative Majorana masses. In ad-

Table 1

Left-handed fermions and scalars.

Particles	$U(1)_D$	$(SU(3)_c, SU(2)_L)_Y$
$Q_i = (u_L \ d_L)_i^T$	0	$(\mathbf{3}, \mathbf{2})_{1/6}$
\bar{u}_{Ri}	0	$(\bar{\mathbf{3}}, \mathbf{1})_{-2/3}$
\bar{d}_{Ri}	0	$(\bar{\mathbf{3}}, \mathbf{1})_{1/3}$
$L_i = (\nu_L \ e_L)_i^T$	0	$(\mathbf{1}, \mathbf{2})_{-1/2}$
\bar{e}_{Ri}	0	$(\mathbf{1}, \mathbf{1})_1$
\bar{Q}_L	−1	$(\mathbf{N}_c, \mathbf{1})_0$
\bar{Q}_R	1	$(\mathbf{N}_c, \mathbf{1})_0$
H	0	$(\mathbf{1}, \mathbf{2})_{1/2}$
η_a	$(-1)^a$	$(\mathbf{N}_c, \mathbf{2})_{1/2}$

dition to the heavy Dirac messenger fermion, we introduce a pair of $SU(2)$ scalar doublets, in order to ensure that at least two neutrino masses are nonzero, as needed to account for neutrino oscillation data [11,12].

2. Bound-state dark matter with Majorana neutrinos

Here we propose a simple realization of the radiative see-saw neutrino mass generation picture, in which a single colored fermion Dirac messenger Q is introduced along with two colored $SU(2)$ doublet scalars η_a ($a = 1, 2$), see Table 1. The heavy messenger Q is charged under $U(1)_D$ and can act as constituent dark matter, made stable thanks to a $U(1)_D$ symmetry.

* Corresponding author.

E-mail addresses: mario.reig@ific.uv.es (M. Reig), restrepo@udea.edu.co (D. Restrepo), valle@ific.uv.es (J.W.F. Valle), oalberto.zapata@udea.edu.co (O. Zapata).
URL: <http://astroparticles.es> (J.W.F. Valle).

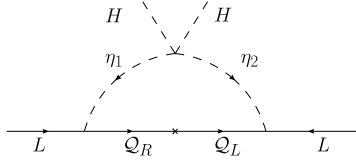


Fig. 1. $U(1)_D$ flux circulating the one-loop neutrino mass.

Note that all dark sector particles carry color. For definiteness, we assign them to the octet $SU(3)_c$ representation. This ensures the viability of the bound-state DM scenario [1]. Notice, however, that our neutrino mass discussion also holds if they had different $SU(3)_c$ transformation properties. The $U(1)_D$ symmetry implies that Q has only a Dirac-type mass term. Hence the new terms in the Lagrangian are the following

$$\mathcal{L} \supset - \left[h_i \overline{Q_R} \tilde{\eta}_1^\dagger L_i + y_i \overline{Q_L} \tilde{\eta}_2^\dagger L_i + M_Q \overline{Q_R} Q_L + \text{h.c.} \right] - \mathcal{V}(H, \eta_a), \quad (1)$$

where summation is implied over repeated indices, and trace over $N_c = 8$ is implicit. The Higgs potential contains the following terms

$$\begin{aligned} \mathcal{V}(H, \eta_a) = & V(H) + V(\eta_a) + \left[\lambda_{3\eta H}^{\dagger} (H^\dagger H) \text{Tr}(\eta_a^\dagger \eta_b) \right. \\ & + \lambda_{4\eta H}^{\dagger} \text{Tr}[(\eta_a^\dagger H)(H^\dagger \eta_b)] + \text{h.c.}] \delta_{ab} \\ & + \text{Tr}[\lambda_{\eta_1 \eta_2} (\eta_1^\dagger \eta_2)(\eta_2^\dagger \eta_1)] \\ & + \left\{ \text{Tr}[\lambda_{\eta_1 \eta_2 H} (H^\dagger \eta_1)(H^\dagger \eta_2)] + \text{h.c.} \right\} \end{aligned} \quad (2)$$

Note that terms like $(\eta_a^\dagger H)^2$ are also forbidden and, as a consequence, the real and imaginary parts of the scalars with nonzero dark charges are degenerate. Moreover, the CP-even and CP-odd scalars do not mix, thanks to our CP conservation assumption.

The Higgs boson is the same as that in the standard model. The neutral scalar components η_a^0 of the dark charge carrying scalars η^a , with quadratic mass coefficients μ_a in Fig. 1, mix into two complex mass eigenstates S_a . Since the tree-level mass term is forbidden by symmetry, neutrino masses are calculable at one-loop order, by the Feynman diagram displayed in Fig. 1.

One finds the following effective mass matrix

$$(\mathcal{M}_\nu)_{ij} = N_c \frac{M_Q}{32\pi^2} (h_i y_j + h_i y_j) \frac{\lambda_{\eta_1 \eta_2 H} v^2}{m_{S_2^0}^2 - m_{S_1^0}^2} \times \left[F\left(\frac{m_{S_2^0}^2}{M_Q^2}\right) - F\left(\frac{m_{S_1^0}^2}{M_Q^2}\right) \right], \quad (3)$$

where $F(m_{S_\beta}^2/M_Q^2) = m_{S_\beta}^2 \log(m_{S_\beta}^2/M_Q^2)/(m_{S_\beta}^2 - M_Q^2)$ 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. Note that the effective one-loop induced neutrino mass matrix has in general rank two, implying two non-vanishing neutrino masses (similar general structure occurs also in non-colored models [13–15]), as required to account for neutrino oscillation data.

As a simple numerical estimate, assuming $\mu_{\eta_1}^2 \gg M_Q^2$, we consider the case $\mu_{\eta_1}^2 = \mu_{\eta_2}^2 \gg \lambda_{\eta_1 \eta_2 H} v^2$ and $\lambda_{3\eta H}, \lambda_{4\eta H} \ll 1$. Taking $m_{S_2}^2 - m_{S_1}^2 = \lambda_{\eta_1 \eta_2 H} v^2$ and $m_{S_{2R}}^2 + m_{S_{1R}}^2 = 2\mu_{\eta_1}^2$, we find

$$(\mathcal{M}_\nu)_{ij} \sim 0.04 \text{ eV} \left(\frac{M_Q}{12.5 \text{ TeV}} \right) \left(\frac{\lambda_{\eta_1 \eta_2 H} v^2}{0.1 \text{ GeV}^2} \right) \left(\frac{15 \text{ TeV}}{\mu_{\eta_1}} \right)^2 \left(\frac{h_i y_j}{10^{-6}} \right). \quad (4)$$

Small neutrino masses are protected by the small parameter $\lambda_{\eta_1 \eta_2 H}$ in whose absence the theory acquires a larger symmetry. As a result, adequate neutrino masses arise naturally for reasonable values of the Yukawa couplings and scalar potential parameters. Due to the flavor structure of the neutrino mass matrix (Eq. (3)), one can express five of the six Yukawa-couplings h_i and y_i in terms of the neutrino observables [16]. Choosing h_1 as a free parameter we have

$$h_i = |h_1| \frac{A_i}{\beta_{11}}, \quad y_i = \frac{1}{2\zeta} \frac{\beta_{ii}}{h_i}, \quad (5)$$

where

$$\beta_{ij} = e^{i\alpha} m_2 V_{i2}^* V_{j2}^* + m_3 V_{i3}^* V_{j3}^*,$$

$$A_j = \pm \sqrt{-e^{i\alpha} m_2 m_3 (V_{12}^* V_{j3}^* - V_{13}^* V_{j2}^*)^2 e^{i2\text{Arg}(h_1)} + \beta_{1j} e^{i\text{Arg}(h_1)}}, \quad \text{for NO,}$$

$$\beta_{ij} = m_1 V_{i1}^* V_{j1}^* + e^{i\alpha} m_2 V_{i2}^* V_{j2}^*,$$

$$A_j = \pm \sqrt{-e^{i\alpha} m_1 m_2 (V_{11}^* V_{j2}^* - V_{12}^* V_{j1}^*)^2 e^{i2\text{Arg}(h_1)} + \beta_{1j} e^{i\text{Arg}(h_1)}}, \quad \text{for IO,}$$

$$\zeta = N_c \frac{M_Q}{32\pi^2} \frac{\lambda_{\eta_1 \eta_2 H} v^2}{m_{S_2^0}^2 - m_{S_1^0}^2} \left[F\left(\frac{m_{S_2^0}^2}{M_Q^2}\right) - F\left(\frac{m_{S_1^0}^2}{M_Q^2}\right) \right].$$

Note that, in the diagonalization condition $U^T \mathcal{M}_\nu U = \text{diag}(m_1, m_2, m_3)$ we have used the PDG form [17] of the lepton mixing matrix [18], namely $U = V P$ and $P = \text{diag}(1, e^{i\alpha/2}, 1)$.

3. Stability of bound-state dark matter and neutrino masses

As suggested in [19] the radiative nature of neutrino mass generation is used to ensure dark matter stability. Moreover, this idea is combined with the proposal that dark matter emerges in the form of stable neutral hadronic thermal relics, as a neutral bound-state of colored constituents, such as QQ , where Q is a vector-like color octet isosinglet fermion (for a general discussion of the cosmology of a stable colored relic see [20]). A necessary and sufficient condition for dark matter stability in this case is the presence of a global $U(1)_D$ dark baryon number,¹ under which the Q is charged [1]. In our present model construction such symmetry also gives rise to radiative neutrino masses. In fact, both dark matter stability, and the Dirac nature of the exotic fermion Q are equivalent, resulting from dark charge conservation.

Note that, 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 Q , and QQ annihilation processes mediated by η are forbidden by the conserved dark symmetry. An adequate thermal relic density of bound-state dark matter requires the lightest constituent vector-like color octet Dirac fermion, Q , to have a mass ≈ 12.5 TeV, so that the QQ hadron weighs approximately 25 TeV [1].

Bound-state dark matter made up by our Dirac octets will be seen in direct searches for nuclear recoil at underground dark matter experiments. The relevant spin-independent cross-section is given as

$$\sigma_{SI} \approx 5.2 \times 10^{-46} \text{ cm}^2 \left(\frac{25 \text{ TeV}}{M_{QQ}} \right)^6 \frac{\Omega_{QQ}}{\Omega_{\text{Planck}}}, \quad (6)$$

¹ Quantum gravity effects may require to gauge global symmetries [21]. An extension where $U(1)_D$ is promoted to a local symmetry can be easily envisaged in a way similar to that considered in [22].

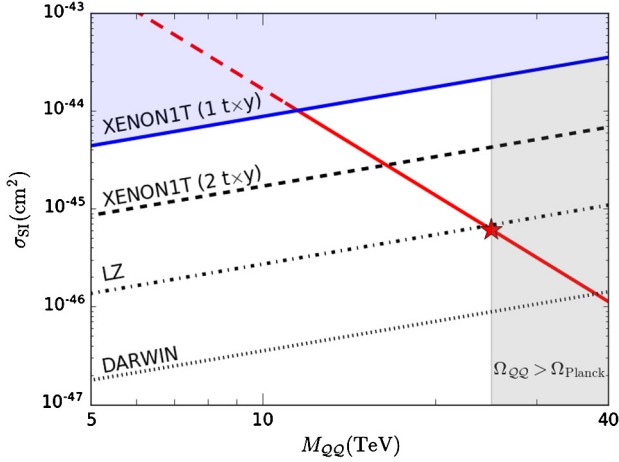


Fig. 2. Spin-independent cross section as a function of $M_{Q\bar{Q}} = 2M_Q$ (red). The star represents the mass required for a thermal bound state 25 TeV dark matter. Lower values can be probed by direct searches, the current bound is indicated in blue, while the black lines (dashed, dotted and dot-dashed) correspond to future sensitivities.

and depends rather sensitively on the dark matter mass $M_{Q\bar{Q}} = 2M_Q$, as shown in the red line in Fig. 2.

Note that the star in the figure assumes that the bound-state DM makes up 100% of the cosmological dark matter. If an additional dark matter component is present, e.g. made of axions, then bound-state dark matter masses below 25 TeV become allowed, as indicated by the red line. In this case their spin-independent cross section would be larger, though their share in the relic density will be lower. The blue line denotes the current XENON1T limit after 1.0 t×yr exposure [23]. This should be compared with the future sensitivities expected at XENON1T [24], LZ [25] and DARWIN [26] indicated by the black (dashed, and dot-dashed) lines. We also note that, within the standard thermal cosmological scenario, DM masses above 25 TeV are ruled out by current observations of the Planck collaboration [27] (gray band). Notice also that the current LHC limit of 2 TeV [28,29], implies that the cross section is always small enough so as to have the bound-state dark matter candidate reaching underground detectors.

4. Neutrinoless double beta decay

In contrast to the proposal in Ref. [9], here total lepton number is a broken symmetry, hence we expect neutrinoless double beta decay to occur. The effective mass parameter characterizing the amplitude for neutrinoless double beta decay is given by [30]

$$\begin{aligned} \langle m_{ee} \rangle &= \left| \sum_j U_{ej}^2 m_j \right| \\ &= \left| c_{12}^2 c_{13}^2 m_1 + s_{12}^2 c_{13}^2 m_2 e^{2i\phi_{12}} + s_{13}^2 m_3 e^{2i\phi_{13}} \right|, \end{aligned} \quad (7)$$

where m_i are the neutrino masses, c_{12} and s_{13} correspond to the angles measured from oscillations and ϕ_{12} , ϕ_{13} are the Majorana phases (here we use the symmetrical parametrization of the lepton mixing matrix [18]).

Since our model predicts the lightest neutrino to be massless, $m_1 = 0$, it follows that there is one single physical (relative) Majorana phase $\phi \equiv \phi_{12} - \phi_{13}$. Furthermore, one can also write the three physical masses directly in terms of the squared mass splittings measured in neutrino oscillation experiments. Depending on the ordering these masses read

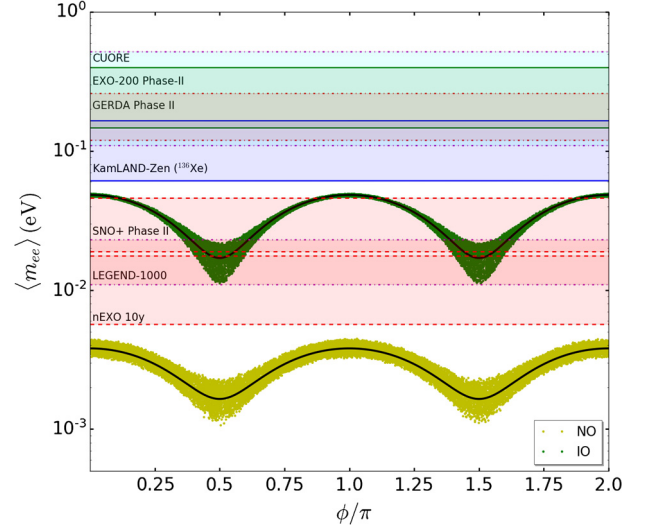


Fig. 3. Effective Majorana mass as a function of the Majorana phase for normal (yellow band) and inverted (green band) mass orderings. The bands represent the 3σ uncertainties in the neutrino oscillation parameters [11,12]. The solid black lines represent their best fit values. The horizontal bands are the experimental limits and sensitivities.

$$\begin{aligned} \text{NO: } m_2 &= \sqrt{\Delta m_{21}^2}, \quad m_3 = \sqrt{\Delta m_{31}^2}, \\ \text{IO: } m_1 &= \sqrt{\Delta m_{13}^2}, \quad m_2 = \sqrt{\Delta m_{13}^2 + \Delta m_{12}^2}. \end{aligned} \quad (8)$$

Given that the lightest neutrino is massless, one can plot the effective Majorana mass parameter $\langle m_{ee} \rangle$ as a function of the unknown Majorana phase ϕ , without loss of generality, see Fig. 3. The results are given for the cases normal (yellow band) and inverted (green band) mass orderings, varying the neutrino oscillation parameters within 3σ [11,12] of their best fit values. One sees that, in contrast to the general three-neutrino scenario, here the $0\nu\beta\beta$ amplitude never vanishes, even when the neutrino mass ordering is of the normal type (models with this feature typically require the existence of specific flavor symmetries [31–33]). The top four horizontal bands represent the 90% C.L. upper limits from CUORE [34], EXO-200 Phase II [35], GERDA Phase II [36] and KamLAND-Zen [37] experiments. The sensitivity bands for the upcoming nEXO experiment after 10 years of data taking [38] as well as for the SNO+ Phase II [39] and LEGEND [40] experiments are indicated by the horizontal red bands.

5. Coupling evolution

Given that we have colored multiplets it is interesting to illustrate their effect in the running of coupling constants. At one loop level the evolution of gauge couplings is governed by

$$\alpha_i(\mu_2)^{-1} = \alpha_i(\mu_1)^{-1} - \frac{b_i}{2\pi} \ln\left(\frac{\mu_2}{\mu_1}\right), \quad (9)$$

where the b_i coefficients are determined by

$$b_i = -\frac{11}{3}C_2(G) + \frac{2}{3} \sum_f T_f d_f + \frac{1}{3} \sum_s T_s d_s. \quad (10)$$

For our model, they are given as

$$\begin{aligned} b_3 &= -7 + 4 \times n_Q + 2 \times n_\eta, \\ b_2 &= -\frac{19}{6} + \frac{4}{3} \times n_\eta, \\ b_1 &= \frac{41}{10} + \frac{4}{5} \times n_\eta, \end{aligned} \quad (11)$$

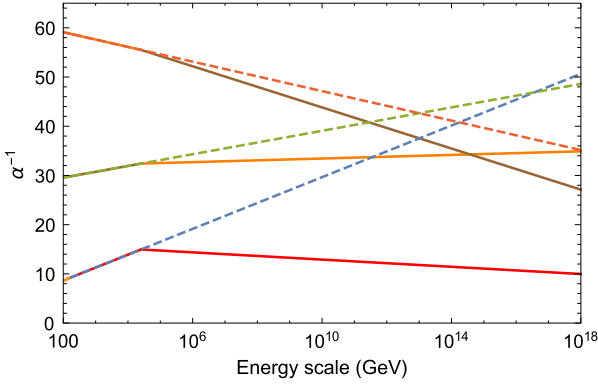


Fig. 4. Running of gauge couplings with light scalars: $M_{\eta_a} \approx 25$ TeV. For DM stability, $M_{\eta_a} > M_Q$ is required.

where n_η is the number of colored scalar fields $\eta \sim (8, 2, 1/2)$. One can easily see that asymptotic freedom can be lost due to the presence of the octets. However, in our case $n_\eta = 2$, hence no Landau pole appears up to the Planck scale, as long as the scalars are heavier than the fermion octets $Q \sim (8, 1, 0, 1)$. This is illustrated in Fig. 4.

Notice that the situation displayed in Fig. 4 seems to threaten the idea of unification of couplings. This, however, is not the full story. In a Grand Unified Theory (GUT) all the fields come in representations of the GUT group. Complete representations of the GUT group do not spoil unification, but change the value of the coupling constant at M_{GUT} . A particular GUT embedding of the color octets $Q \sim (8, 1, 0)$, however, lies beyond the scope of this work.

6. Phenomenology

The cosmological relic abundance requires $M_Q \approx 12.5$ TeV, which lies beyond the kinematical reach of the LHC [1]. A future hadron collider with a center of mass energy of 100 TeV, would probe masses up to $M_Q \lesssim 15$ TeV [41]. Hence, given the phase space, the Q 's would be copiously pair produced through the processes $q_i \bar{q}_i \rightarrow Q \bar{Q}$ and $gg \rightarrow Q \bar{Q}$, and would be expected to hadronize. Different bound states such as neutral hybrid states Qg (detected as neutral hadrons, presumably stable) or charged $Qq\bar{q}'$ states, or more exotic $Qqqq$ states are expected to form after on-shell Q production. They would give rise to visible signals at such future hadron colliders, making them a unique environment to test colored dark matter.²

In contrast with bound-state dark matter with radiative Dirac neutrino masses induced by color octets circulating the loop [9], in the Majorana case the color octet scalars can lie at the same scale as the color octet fermion. In fact, if the lightest color octet scalar, η_1 , transforming as a weak doublet, has a mass close to that of Q , they could be pair-produced with similar cross sections.³ The pair produced scalars further decay into a $Q\bar{Q}$ pair and charged leptons or neutrinos. This gives rise to similar signals with long-lived bound states but with extra charged leptons or missing transverse energy.

Note also that our bound-state dark matter constituents can mediate charged lepton flavor violation effects, whose rates will depend on the scalar masses and may reach detectability levels.

² Concerning the scalar colored particles, they may also mediate lepton number violating processes at colliders. However, since they lie at scales of 12.5 TeV or higher, the corresponding cross section is expected to be very small.

³ We compared the reported cross sections of [41] for the color octet Dirac fermion with the MadGraph [42] output of FeynRules [43] for color-octet isospin-doublet scalar implemented by [44].

7. Summary and outlook

We have proposed a simple and viable theory in which dark matter emerges as a stable neutral hadronic thermal relics, whose stability results from an exact $U(1)_D$ symmetry. Neutrinos pick up radiatively induced Majorana masses from the exchange of colored dark matter constituents, giving a common origin for both dark matter and neutrino mass, with a lower bound for neutrinoless double beta decay and direct tests at direct DM searches at nuclear recoil experiments. The scheme provides a consistent ultraviolet complete setup, free of Landau poles all the way up to the Planck scale, provided the scalars are heavy enough. The new states may also lead to other phenomenological signals at future hadron colliders as well as lepton flavor violation experiments.

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