Is it possible to explain neutrino masses with scalar dark matter?

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We present a scenario in which a remarkably simple relation linking dark matter properties and neutrino masses naturally emerges. This framework points towards a low energy theory where the neutrino mass originates from the existence of a light scalar dark matter particle in the keV-MeV mass range. We discuss different ways to constrain and test this scenario by means of astrophysical and cosmological observations as well as laboratory experiments. Finally, we point out that one interesting aspect is that the implied mass range is compatible with the one required for the explanation of the mysterious emission of 511 keV photons from the center of our galaxy in terms of dark matter annihilation into e^+e^- pairs.

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I. INTRODUCTION

The discovery of nonzero neutrino masses in neutrino oscillation experiments [1] and the increasing evidence for about 23% of the content of the Universe being in the form of dark matter [2] are the two main indications for physics beyond the standard model. These two issues, the origin of neutrino masses and the nature of dark matter, have been long-standing problems in particle physics. Yet, in general, they are considered as two different topics and current explanations rely on completely different mechanisms, involving unrelated particles and scales. Although there have been some proposals to establish a link [3–5], a simple and natural picture in which the neutrino mass scale, the dark matter properties, and dark matter abundance would be quantitatively related is still missing. In particular, to the best of our knowledge, there is no model in the literature which uniquely determines the dark matter scale and predicts, at the same time, a direct connection between the smallness of neutrino masses and the observed dark matter relic density.

In this paper, we present a scenario where such a prediction exists and therefore establish a quantitative link between these two fields. For this purpose, we extend the standard model Lagrangian with the following interaction term so as to describe the existence of a neutral (singlet under $SU(2)_L \times U(1)$) scalar particle ϕ uniquely (or predominantly) coupled to neutrinos:

$$\mathcal{L}_I \supset g\phi\bar{N}\nu. \tag{1}$$

Here g is a coupling constant, N is a Majorana neutrino (with a mass m_N), ϕ plays the role of dark matter (hereafter referred to as the SLIM particle for Scalar as LIght as MeV), and ν is the standard left-handed neutrino. Since the mass of the particle N is of Majorana type, lepton number is not conserved. The term $\mathcal{L}_I$ under consideration explicitly breaks the electroweak symmetry and should be therefore regarded as an effective Lagrangian.

In this paper, we assume that SLIMs interact only with neutrinos. However, one could extend this scenario so as to include SLIM interactions with electrons [6–9] or mixing effects [10] which could induce signatures in future neutrino experiments. Although there is a rich amount of literature about scalar dark matter particles with significant interactions with neutrinos (see e.g. Refs. [11,12]), the link between neutrino mass and the dark matter density in this scenario has been overlooked.

As we will see, in this scenario not only a remarkably simple relationship between the dark matter cross section and the neutrino mass scale naturally emerges, but the requirement of sub-eV neutrino masses, as imposed by experimental constraints, points towards light dark matter particles (with a mass of a few MeV or less). Our expression therefore suggests that the issues regarding the dark matter and neutrino masses could be not only closely related, but they could also share the same low energy origin. This may be the first step towards a low energy theory beyond the standard model.

II. LINKING DARK MATTER AND NEUTRINO MASS

In the Lagrangian given in Eq. (1), the SLIM particle, i.e., the scalar ϕ, can be either real or complex, and we consider the case when it is lighter than the Majorana neutrino, $m_N > m_\phi$. In this case, the particle ϕ is stable and constitutes our dark matter candidate. In contrast N decays into ϕ and ν with a decay rate $\Gamma_N = g^2 m_N^2/(16\pi E_N)$. For simplicity, in this section we drop the flavor indices, although they will be revisited in Sec. V.

From the Lagrangian $\mathcal{L}_I$, left-handed neutrinos acquire a mass term m_ν via the one-loop correction depicted in Fig. 1 [11]. In the case of a real scalar field ϕ, this contribution is given by

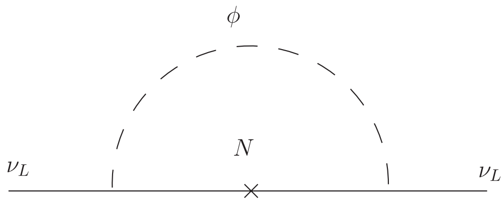

FIG. 1. Left-handed neutrinos acquire a very small mass due to their interactions with a SLIM and a N particle.

$$m_\nu = \frac{g^2}{16\pi^2} m_N \left[\ln\left(\frac{\Lambda^2}{m_N^2}\right) - \frac{m_\phi^2}{m_N^2 - m_\phi^2} \ln\left(\frac{m_N^2}{m_\phi^2}\right) \right], \quad (2)$$

with g, m_N, m_ϕ, and Λ (the ultraviolet cutoff of the effective theory) being free parameters.

This mechanism is the same as that in Refs. [4,5,11] except from the fact that, in our scenario, ϕ is a singlet under the electroweak symmetry. Like in Refs. [4,5], we assume that ϕ does not acquire a vacuum expectation value (VEV), so Eq. (1) does not induce any tree-level contribution to the left-handed neutrino mass. In this scenario, light neutrinos are Majorana particles (see Ref. [13] for a proof within the context of supersymmetry). This could be tested in neutrinoless double beta decay experiments [14].

If SLIM interactions with standard model particles are limited to Eq. (1), their annihilation in the early Universe will be given by the three diagrams depicted in Fig. 2, corresponding to annihilations into (anti)neutrinos pairs. The sum of these three contributions sets the annihilation rate and therefore determines the SLIM relic density.

The SLIM pair annihilation cross sections into neutrino $\sigma(\phi\phi \to \nu\nu)$ and antineutrino $\sigma(\phi\phi \to \bar{\nu}\,\bar{\nu})$ pairs (see first two diagrams of Fig. 2) are equal. The cross section into neutrino-antineutrino pairs (last diagram of Fig. 2) is suppressed with respect to the other two by a factor m_ν^2/m_N^2.[1] Hence the total annihilation cross section times v_r (the relative velocity of the initial state particles) approximately corresponds to $2\langle\sigma(\phi\phi \to \nu\nu)\mathrm{v}_r\rangle$ with

$$\langle\sigma(\phi\phi \to \nu\nu)\mathrm{v}_r\rangle = \langle\sigma(\phi\phi \to \bar{\nu}\,\bar{\nu})\mathrm{v}_r\rangle \simeq \frac{g^4}{4\pi} \frac{m_N^2}{(m_\phi^2 + m_N^2)^2}, \quad (3)$$

where the notation $\langle \ldots \rangle$ denotes the thermal average.

The interesting point is that, for $m_\phi < m_N \ll \Lambda$, Eq. (2) can be rewritten in a very convenient way using Eq. (3) as

$$m_\nu \simeq \sqrt{\frac{\langle\sigma\nu_r\rangle}{128\pi^3}} m_N^2 \left(1 + \frac{m_\phi^2}{m_N^2}\right) \ln\left(\frac{\Lambda^2}{m_N^2}\right). \quad (4)$$

To fit the observed abundance, the total SLIM pair annihi-

[1] In the case of complex particles, the cross section into neutrino-antineutrino is suppressed with respect to the first two by only a factor $p_{\rm dm}^2/m_N^2$, where $p_{\rm dm}$ is the dark matter momentum.

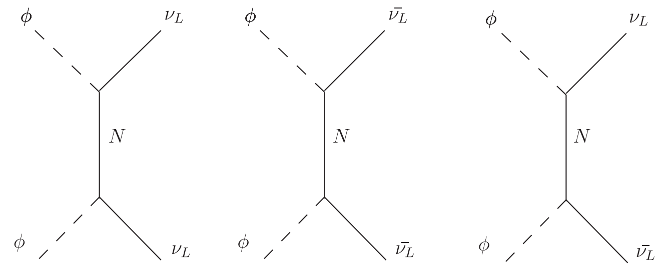

FIG. 2. Here are the three diagrams contributing to the tree-level total SLIM pair annihilation cross section.

lation cross section must be of the order of [15]

$$\sigma\nu_{r|_{\rm ref}} \simeq 7\,10^{-27} \frac{x_F}{\sqrt{g_\star}} \left(\frac{\Omega_{\rm dm} h^2}{0.1}\right)^{-1} \mathrm{cm}^3/\mathrm{s}, \quad (5)$$

where $\Omega_{\rm dm}$ is the cosmological parameter associated with dark matter, h is the normalized Hubble constant, $x_F \in [12,16]$ for particles in the MeV-100 GeV range and $g_\star$ is the number of relativistic degrees of freedom at the dark matter chemical decoupling. Equation (5) is therefore almost independent of the dark matter mass and about 10^{-26} cm^3/s. Inserting $\sigma\nu_{r|_{\rm ref}}$ into Eq. (4), we now obtain

$$\left(\frac{m_\nu}{10^{-2}\ \mathrm{eV}}\right) \simeq 0.039\,\mathcal{A}\,(\Omega_{\rm dm}h^2)^{-(1/2)} \left(\frac{m_N}{\mathrm{MeV}}\right)^2, \quad (6)$$

with $\mathcal{A} = \sqrt{x_F} g_\star^{-1/4} (1 + \frac{m_\phi^2}{m_N^2}) \ln(\frac{\Lambda^2}{m_N^2})$. This relation shows that, if dark matter is a scalar which does not acquire a VEV and predominantly annihilates into Majorana neutrinos, the left-handed neutrino masses and the dark matter abundance *are very strongly related*.

The quantities m_ν and $\Omega_{dm}h^2$ are measured by neutrino experiments and cosmological observations such as the cosmic microwave background (CMB) experiments, respectively. Hence the only free parameters in Eq. (6) are m_N, m_ϕ/m_N, and Λ. Since the dependence of Eq. (4) on Λ is only logarithmic and since varying the ratio m_ϕ/m_N between 0 and 1 does not modify the result of Eq. (4) by more than a factor 2, m_N is the only important free parameter. Equation (6) is therefore extremely predictive.

Let us take, for instance, $\Lambda \sim E_{\rm electroweak} \sim 200$ GeV and consider the experimental constraints $0.05\ \mathrm{eV} < m_\nu < 1\ \mathrm{eV}$, we then obtain:

$$O(1)\ \mathrm{MeV} \lesssim m_N \lesssim 10\ \mathrm{MeV}. \quad (7)$$

Equation (7) is remarkable. Not only does it indicate the existence of an *upper* limit on the dark matter mass but it also shows that the SLIM mass does not correspond to the high energy scale that is usually considered. In addition, it means that N is an electroweak singlet or has very weak couplings to the standard model Z boson. Note that this range will be narrowed down by improving the bounds on neutrino masses or possibly, by directly measuring them.

Since the coupling g (that enters Eq. (3)) is given by

$$g \simeq 10^{-3}\sqrt{\frac{m_N}{10 \text{ MeV}}}\left(\frac{\langle\sigma v_r\rangle}{10^{-26} \text{ cm}^3/\text{s}}\right)^{1/4}\left(1+\frac{m_\phi^2}{m_N^2}\right)^{1/2}, \quad (8)$$

one obtains by inserting Eq. (7) into Eq. (8):

$$3\times 10^{-4} \lesssim g \lesssim 10^{-3}. \quad (9)$$

To be accurate, one should take into account flavor effects, i.e., one should specify the combination of neutrino flavors to which N is coupled, keeping in mind that at least a second (heavier) N is necessary to lead to at least two massive neutrinos. This would be done in the next section. Note, however, that the main conclusions of this section remain unchanged.

It is well known that hot dark matter generates a damping of primordial fluctuations at the scales of galaxy clusters. Hence strong constraints on its mass and energy density can be obtained [17]. On the other hand, warm dark matter, with a smaller free-streaming (or collisional damping) length, does not have such a fatal implication. It can even solve some of the problems of cold dark matter scenarios [18]. Large scale structure arguments impose that the mass of a warm dark matter candidate should be greater than a few keV [19]. Hence we can conclude that m_ϕ lies between

$$\text{a few keV} < m_\phi < 10 \text{ MeV}.$$

Let us now discuss the case of a complex scalar field, $\phi = (\phi_1 + i\phi_2)/\sqrt{2}$ where ϕ_1 and ϕ_2 are real fields with masses m_{ϕ_1} and m_{ϕ_2}. The various equations obtained for real ϕ are modified but the overall picture remains the same. In particular, Eq. (2) becomes

$$m_\nu = \frac{g^2}{32\pi^2} m_N\left[\frac{m_{\phi_1}^2}{(m_N^2 - m_{\phi_1}^2)}\ln\left(\frac{m_N^2}{m_{\phi_1}^2}\right) - \frac{m_{\phi_2}^2}{(m_N^2 - m_{\phi_2}^2)}\ln\left(\frac{m_N^2}{m_{\phi_2}^2}\right)\right], \quad (10)$$

and Eq. (4) changes into

$$m_\nu = \sqrt{\frac{\langle\sigma v_r\rangle}{128\pi^3}}\left[m_{\phi_1}^2\ln\left(\frac{m_N^2}{m_{\phi_1}^2}\right) - m_{\phi_2}^2\ln\left(\frac{m_N^2}{m_{\phi_2}^2}\right)\right], \quad (11)$$

where we have neglected the terms suppressed by $m_{\phi_{1,2}}^2/m_N^2$. Note that the cutoff dependence drops out in Eq. (10). This comes from the minus sign between the two ϕ component contributions. This negative sign is due to the fact that if both states have the same mass, lepton number is conserved (assigning $L = -1$ to ϕ and $L = 0$ to N) and the neutrino mass must therefore vanish. This is different from the real ϕ case where lepton number is necessarily violated if N has a Majorana mass.

In Eq. (11), for small $\phi_1 - \phi_2$ mass splittings, the neutrino mass is determined by the quantity $m_{\phi_1}^2 - m_{\phi_2}^2$ while, in Eq. (4), it was determined by m_N^2. Hence, instead of Eq. (7), we now obtain

$$(1 \text{ MeV})^2 \lesssim |m_{\phi_1}^2 - m_{\phi_2}^2| \lesssim (20 \text{ MeV})^2. \quad (12)$$

For definiteness, we have assumed that the ratio m_N/m_{ϕ_1} was ranging from 10 to 10^5. In Eq. (11), the mass m_N is a free parameter which can be much larger than the mass of the Z boson. Hence, in the complex case (unlike the real case), the particle N can have electroweak couplings.

If, for instance $m_{\phi_2} < m_{\phi_1}$, one would expect the unstable particle, ϕ_1, to decay into ϕ_2 plus a neutrino-antineutrino pair. The ϕ_2 particle, being stable, would be our dark matter candidate. Note also that if the mass splitting between m_{ϕ_1} and m_{ϕ_2} is small, one has to take into account the coannihilations between ϕ_1 and ϕ_2 for the calculation of the dark matter relic density. This may slightly change Eqs. (11) and (12) but this detailed study is beyond the scope of the present paper.

In summary, if ϕ is a real field, the natural scale for the dark matter mass is the MeV range or below. If ϕ is a complex field, a suitable scale is also the MeV range, although Eq. (12) does not uniquely predict that the dark matter mass must lie in the low energy range.

Notice that if m_N is large, above few GeV, as it might be in the complex case, the self-energy correction to the ϕ mass which is induced by the g interaction is larger than the tree-level mass itself. This is similar to the Higgs mass correction induced by the Yukawa couplings in the usual seesaw model, which is larger than the Higgs mass for $m_N > 10^7$ GeV. For m_N as large as 100 GeV and $m_\phi = 10$ MeV, the mass stability requires a fine-tuning of order one per few thousands. This fine-tuning could be reduced if we assume a larger m_ϕ and a smaller ϕ mass splitting. This issue might be solved in the more general context of theories which embed the model we propose.

All in all, obtaining the MeV scale is quite an amazing finding since this corresponds to the dark matter mass range which, for instance, is also required to explain the 511 keV emission line from the center of our galaxy [6–8]. To this aim, one needs to introduce an additional interaction between the SLIM and electrons so that the SLIM particles can pair annihilate and produce positrons.

III. CAN SLIMS BE CONSTRAINED?

The scenario that we discuss in this paper is not excluded by the bounds from direct and indirect dark matter detection experiments. It is interesting to note that future neutrino detectors will be able to directly test this scenario if the dark matter mass is above ~10 MeV [20]. In addition, SLIM particles are consistent with the constraints from large scale structure formation and, in particular, with the constraints obtained in Ref. [21].

We now discuss a set of constraints and possible tests concerning light scalars or neutrino-scalar interactions that might be relevant in our case.

A. Nucleosynthesis

SLIM particles could affect primordial big bang nucleosynthesis (BBN) if they were too light. As shown in Ref. [22], for masses much smaller than 1 MeV, each scalar contribution to the primordial element abundances is equivalent to $4/7$ new degrees of freedom (dof) and each fermion to 1 dof, which would increase the helium and deuterium abundance, and decrease the lithium one. If one only considers helium, the results of Ref. [22] strongly disfavor values of the mass below 1 MeV. However, it is a conservative bound given that the combination of the CMB determination of the baryon-to-photon ratio with primordial light-element abundances observations or with large scale structure data, allows up to $\sim$1.5 extra dof (at 95% CL) [23].

For masses above $\sim$10 MeV, there is no effect on BBN due to the fact that either the dark matter density is Boltzmann suppressed or dark matter has completely annihilated before the neutrino-electron decoupling. For intermediate masses (between $\sim$1 and 10 MeV), the impact of the dark matter energy density on BBN is negligible (due to the Boltzmann suppression factor), but the secondary neutrino production could affect the relation $T_\nu(T)$. In the case of a scalar dof, the authors of Ref. [22] showed that this is, for the case of the helium (deuterium and lithium) abundance, equivalent to adding less than 0.5 new dof for SLIM masses down to 1 MeV (2–4 MeV). Our result combined with the analysis in Refs. [22,23] therefore suggests that, in the real case, the SLIM mass must range between $\sim$1 MeV to at most 10 MeV while in the complex case it should just be larger than $\sim$1 MeV.

B. Cosmic microwave background

In Ref. [24], the authors have derived strong bounds on the coupling of the neutrino to a light scalar particle from the CMB. The strongest bound on the coupling comes from invisible decays such as $\nu \to \nu' \phi$ (with $m_\phi < 0.1$ eV). Such decays cannot occur in our case since the SLIM particles are much heavier than 0.1 eV. Moreover, in Ref. [24], they also constrain the coupling by studying "binary interactions" such as $\nu\bar{\nu} \to \phi\phi$, $\nu\nu \to \nu\nu$, and $\nu\phi \to \phi\nu$. However, these processes are not relevant for the case presented here. At $T \sim 0.1$ eV indeed, $\nu\bar{\nu} \to \phi\phi$ ($\nu\nu \to \phi\phi$) is not possible because $m_\phi \gtrsim$ MeV. The process $\nu\nu \to \nu\nu$ could happen through a box diagram but it is suppressed by a factor of $g^8/(16\pi^2)^2$. At last, one could consider $\nu\phi \to \nu\phi$. However, as shown in Ref. [6], at the recombination epoch, the neutrino-dark matter elastic scattering cross section vanishes at the local limit ($m_N m_\phi \gg T$).[2]

[2]At $T \sim 0.1$ eV, the $\nu\phi \to \bar{\nu}\phi$ cross section is given by the square of the neutrino temperature and is therefore extremely small.

C. Structure formation

Similar to photon-dark matter interactions [25], a large neutrino-dark matter elastic scattering cross section at the recombination epoch could in principle affect structure formation [16]. However the constraint derived in Ref. [16] does not apply in our case. As explained previously, at the recombination epoch the associated cross section strictly vanishes in the local limit or is extremely suppressed [6]. Hence, as predicted in Ref. [6], our model does not change structure formation on visible scales.

D. Core-collapse Supernovae

Large neutrino-SLIM coupling can also be constrained by studying the core-collapse Supernovae, as pointed out in Ref. [26]. In Ref. [26], a model is investigated in which light dark matter is coupled to a new gauge boson resulting in cross sections for the elastic interactions of the electron and neutrino with dark matter of the same order of magnitude as the dark matter pair annihilation cross section into electron-positron and neutrino pairs. More importantly, they suppose that the dark matter coupling to nucleons is large enough to lead to a nucleon-dark matter elastic scattering cross section ranging from 10^{-40} cm^2 up to 10^{-35} cm^2. With such interactions to nucleons, one obtains a dark matter-sphere temperature[3] that is smaller than $\sim$3 MeV and therefore changes the neutrino-sphere temperature. Notice, however, that our model here is different. In our scenario, the SLIMs interact only with neutrinos. As a result, the neutrino-sphere temperature remains the same. Hence, our very Lagrangian cannot be constrained from the modification of the neutrino-sphere temperature.

Even if we extend our Lagrangian to include dark matter interactions with nucleons, we will end up with a dark matter-nucleon elastic scattering cross section that is smaller than 10^{-40} cm^2. The SLIM couplings to nucleons could be larger than the coupling to neutrinos but the mass of the charged particle that is associated with this interaction, m_{F_q}, must be much larger than m_N to satisfy the LEP and Tevatron constraints, with a cross section that goes as $T^6/m_{F_q}^8$ (assuming a scalar coupling). Hence, even if we introduce interactions with nucleons, the SLIM-nucleon cross section would be very much suppressed with respect to the cross section considered in Ref. [26]. Hence the conclusions of Ref. [26] do not apply to our case.

On the other hand, large neutrino-dark matter interactions could modify the neutrino diffusion scale and change the cooling time of the supernova. We estimate the corresponding cross section in our model to lie between $[10^{-40}, 10^{-38}]$ cm^2 for a neutrino temperature of 30 MeV and $1 < m_N < 10$ MeV. Taking a dark matter number density $n_{eq,X}$ equal to $n_N/100$ (for T = 30 MeV) and using

[3]We use the same definition as in Ref. [26] for the dark matter-sphere temperature.

a neutrino-dark matter cross section of about 10^{-36} cm^2, the authors in Ref. [26] found a neutrino diffusion scale of $\lambda = c/(\sigma v|_{\nu-\phi} n_\phi) \sim 0.3$ cm (it is about $\lambda = c/(\sigma v|_{\nu-\mathrm{nuc}} n_{\mathrm{nuc}}) \sim 35$ cm in the standard case). With our cross section, we estimate this diffusion length to be about 30 cm (it can be exactly that of the standard case when the neutrino-SLIM interaction rate is smaller than in the standard case). Therefore SLIM particles are not constrained by SN observations.

It is also interesting to note that, in the case of future supernova neutrino observations, one may be able to test this scenario by studying the neutrino energy spectra. For example, in a similar way as to what happens in the model of Ref. [27], one would expect dips in the diffuse Supernova neutrino spectrum for $m_\phi \sim 10 - 100$ keV [28].

E. Meson decay

As far as laboratory constraints are concerned, light scalar emission has not been observed in pion and kaon decays. For kinematically allowed decays a very conservative bound can be obtained, which constrains the coupling in Eq. (1) to be $g \lesssim 10^{-2}$ [29]. Improving the present experimental bounds seems nevertheless feasible. For real ϕ, the upper bounds on m_ϕ and m_N together with the relatively large value of the coupling (see Eq. (9)) promise *observable effects* in kaon and pion decay experiments. This would make this scenario even more appealing as it could be tested soon.

Many other constraints were discussed in Ref. [6] with the conclusions that a scenario like the one presented here is perfectly viable.

F. Positron emission from the galactic center and 511 keV

In our model there is no interaction with electrons. However, one could extend the model to add a coupling to electrons or to any other standard model particles. One of the conditions that one has to respect is that the dark matter pair annihilation cross section into $e^- e^+$ must be suppressed by at least 4 orders of magnitude with respect to the annihilation cross section into neutrino pairs, to avoid an overproduction of low energy gamma rays and be simultaneously compatible with the 511 keV observation. In addition, the dark matter mass is also constrained to be smaller than a few MeV to avoid an excess of gamma rays due to high order processes via internal bremsstrahlung [30–32] and positron annihilations in flight with the electrons in the interstellar medium [32,33]. On the other hand, the couplings with leptons are also bounded by other astrophysical and laboratory bounds [22,26,34]. Interestingly, a number of different tests of this scenario have been recently proposed [35].

This extended scenario could be achieved, as proposed in Refs. [7,9], through the exchange of a charged particle F_e [6,9] with $m_{F_e} \ggg m_N$. In this case the product of the F_e couplings (to left, c_l, and right, c_r, handed electrons) has to be $c_l c_r \leqslant 1.6710^{-4} (\frac{m_{F_e}}{100\ \mathrm{GeV}})(\frac{m_\phi}{\mathrm{MeV}})$ [9]. To make a comparison with supersymmetry, our SLIM particle would then be somehow similar to a light sneutrino with negligible couplings to the Z (more or less like in Ref. [36]), F_e would be equivalent to a "heavy" chargino and N a light neutralino (but N and F_e would not be coupled to any other SM species except to the neutrino and the electron, respectively).

IV. EXAMPLES OF THEORETICAL REALIZATIONS OF OUR MODEL

There are certainly many ways to obtain the effective low energy ($SU(2)_L \times U(1)$ breaking) term of Eq. (1) from an underlying theory. If ϕ and N are both $SU(2)_L$ singlets, this term is necessarily effective. It can be obtained, after spontaneous electroweak symmetry breaking, from the exchange of an additional scalar doublet [37] or an extra vectorlike fermion singlet (or doublet). The extra particles then have to be well above the MeV scale. If N is not a $SU(2)_L$ singlet (which, as discussed, can be realized only for a complex ϕ), there are various possibilities to obtain the interaction in Eq. (1). For example, if, as in Ref. [6], N is a "mirror" fermion (which along with a charged lepton E_R forms a $SU(2)_L$ doublet), the coupling in Eq. (1) can be considered a "fundamental" one. In Ref. [6], N was not a Majorana particle, so it could not lead to the mechanism described in this paper. However, one can consider a more sophisticated model where N is still in an $SU(2)_L$ doublet but with an "effective" Majorana mass that is induced from spontaneous $SU(2)_L$ symmetry breaking. This mass can be obtained from a mirror seesaw mechanism between N and an extra left-handed $SU(2)_L$ singlet N_L fermion (to be added to the Lagrangian of Ref. [6]). If the mass of the N_L is not far above the electroweak scale, and if the allowed "$N - N_L - H$" Yukawa coupling is large enough, this seesaw can lead to m_N well above ~50 GeV, which can satisfy various collider constraints (e.g., invisible decay width of the Z boson, direct searches for these particles, electroweak precision measurements). For a complex ϕ, such a value of m_N satisfies the requirements of this scenario. In this model, the lightest ϕ component would be stable if, for example, like N and N_L, it is odd under a Z_2 symmetry (similar to models considered in Refs. [4,11]). This solution is particularly interesting since, as shown in Ref. [9], mirror fermions can be at the origin of the 511 keV emission. A systematic study of all these possibilities and related constraints is beyond the scope of this paper.

V. FLAVOR EFFECTS

In this section, we restore the flavor indices in formulae for the neutrino masses and the annihilation rate of the

SLIM particles. The interaction term is

$$\mathcal{L} = g_{i\alpha}\phi \bar{N}_i \nu_\alpha.$$

Restoring the flavor indices in Eq. (2), for real ϕ we find

$$(m_\nu)_{\alpha\beta} = \sum_i \frac{g_{i\alpha}g_{i\beta}}{16\pi^2} m_{N_i}\left(\log\frac{\Lambda^2}{m_{N_i}^2} - \frac{m_\phi^2}{m_{N_i}^2 - m_\phi^2}\log\frac{m_{N_i}^2}{m_\phi^2}\right). \tag{13}$$

Similarly for the complex case, we find (see Eq. (10))

$$(m_\nu)_{\alpha\beta} = \sum_i \frac{g_{i\alpha}g_{i\beta}}{32\pi^2} m_{N_i}\left[\frac{m_{\phi_1}^2}{(m_{N_i}^2 - m_{\phi_1}^2)}\ln\left(\frac{m_{N_i}^2}{m_{\phi_1}^2}\right) - \frac{m_{\phi_2}^2}{(m_{N_i}^2 - m_{\phi_2}^2)}\ln\left(\frac{m_{N_i}^2}{m_{\phi_2}^2}\right)\right]. \tag{14}$$

With only one right-handed neutrino, there will be only three free parameters entering the mass formula. Thus, similar to the canonical seesaw mechanism we will need more than one right-handed neutrino to fit the data. In general, we can have "n" extra right-handed neutrinos contributing to the masses. The neutrino mass matrix can be written as

$$m_\nu = U \cdot \mathrm{Diag}[m_1, m_2 e^{2i\phi_2}, m_3 e^{2i\phi_3}]U^T.$$

It is straightforward to show that, for $n \geqslant 3$, to yield neutrino masses compatible with the data, the coupling matrix has to have the following structure regardless of whether ϕ is real or complex:

$$g = \mathrm{Diag}(X_1, \ldots, X_n)\cdot O \cdot \mathrm{Diag}(\sqrt{m_1}, \sqrt{m_2}e^{i\phi_2}, \sqrt{m_3}e^{i\phi_3})U^T, \tag{15}$$

where O is an arbitrary $n \times 3$ matrix that satisfies $O^T \cdot O = \mathrm{Diag}(1, 1, 1)$. Notice that for $n = 3$, this means O is an arbitrary orthogonal matrix. In the case of real ϕ,

$$X_i = 4\pi\left(\frac{1}{m_{N_i}}\right)^{1/2}\left(\log\frac{\Lambda^2}{m_{N_i}^2} - \frac{m_\phi^2}{m_{N_i}^2 - m_\phi^2}\log\frac{m_{N_i}^2}{m_\phi^2}\right)^{-1/2}, \tag{16}$$

while for complex ϕ,

$$X_i = \sqrt{\frac{32\pi^2}{m_{N_i}}}\left(\frac{m_{\phi_1}^2}{m_{N_i}^2 - m_{\phi_1}^2}\log\frac{m_{N_i}^2}{m_{\phi_1}^2} - \frac{m_{\phi_2}^2}{m_{N_i}^2 - m_{\phi_2}^2}\log\frac{m_{N_i}^2}{m_{\phi_2}^2}\right)^{-1/2}. \tag{17}$$

Now, let us turn our attention to the annihilation cross section. In the case of the real ϕ, restoring the flavor indices in Eq. (3) we find

$$\langle\sigma(\phi\phi \to \nu_\alpha\nu_\beta)v_r\rangle = \langle\sigma(\phi\phi \to \bar{\nu}_\alpha\bar{\nu}_\beta)v_r\rangle = \frac{1}{4\pi}\left|\sum_i \frac{g_{i\alpha}g_{i\beta}m_{N_i}}{m_\phi^2 + m_{N_i}^2}\right|^2. \tag{18}$$

Hence, the total cross section will be

$$\langle\sigma v_r\rangle|_{\mathrm{total}} = \frac{2}{4\pi}\left(\sum_\alpha\left|\sum_i \frac{g_{i\alpha}^2 m_{N_i}}{m_\phi^2 + m_{N_i}^2}\right|^2 + \frac{1}{2}\sum_{\alpha\neq\beta}\left|\sum_i \frac{g_{i\alpha}g_{i\beta}m_{N_i}}{m_\phi^2 + m_{N_i}^2}\right|^2\right) = \frac{1}{4\pi}\left(\sum_\alpha\left|\sum_i \frac{g_{i\alpha}^2 m_{N_i}}{m_\phi^2 + m_{N_i}^2}\right|^2 + \sum_{\alpha,\beta}\left|\sum_i \frac{g_{i\alpha}g_{i\beta}m_{N_i}}{m_\phi^2 + m_{N_i}^2}\right|^2\right). \tag{19}$$

In the case of complex ϕ with $m_2 < m_1$ (ϕ_2 constituting the dark matter), we should replace m_ϕ with m_{ϕ_2} and g with $g/\sqrt{2}$.

Plugging the couplings given in Eq. (15) into Eq. (19), it is easy to check that the lighter m_{N_i} that dominate the annihilation cross section must in general take values similar to the one obtained in the single flavor case. As an example, let us consider a simple model with only two right-handed neutrinos: N_1, N_2. This is the most economic scenario compatible with the present neutrino data. If there are only two-right-handed neutrinos, we will obtain $\mathrm{Det}(m_\nu) = 0$ which means the lightest neutrino eigenvalue vanishes and therefore the hierarchical scheme is the only possibility. Within the normal hierarchical scheme and the scenario with real ϕ, taking $m_{N_1} = m_{N_2}$ and neglecting the corrections of order of s_{13} and $\Delta m_{\mathrm{sol}}^2/\Delta m_{\mathrm{atm}}^2$, where $s_{13} = \sin\theta_{13}$ with θ_{13} being the small mixing angle as written in the standard form of the PMNS matrix and $\Delta m_{\mathrm{sol}}^2$ and $\Delta m_{\mathrm{atm}}^2$ being the solar and atmospheric mass squared difference, respectively. We can write

$$\langle\sigma v_r\rangle|_{\mathrm{total}} = \frac{(4\pi)^3}{(m_\phi^2 + m_N^2)^2}\left(\log\frac{\Lambda^2}{m_N^2} - \frac{m_\phi^2}{m_N^2 - m_\phi^2}\log\frac{m_N^2}{m_\phi^2}\right)^{-2} \times (\Delta m_{\mathrm{atm}}^2)(1 + c_{23}^4 + s_{23}^4), \tag{20}$$

where $c_{23} = \cos\theta_{23}$ and $s_{23} = \sin\theta_{23}$ with θ_{23} being the so-called atmospheric mixing angle. This equation shows that in this case we recover the relation we have found in the case of single flavor analysis, up to a factor of order unity. Notice that in the final result, the dependence on the arbitrary matrix O disappears (see Eq. (15)). In the general

TABLE I. Some possible solutions in the framework of real ϕ. "N" and "I," respectively, denote normal and inverted mass scheme. We have taken $\langle\sigma v_r\rangle = 10^{-26}$ cm^3/s, $\Delta m_{\mathrm{atm}}^2 = 2.6 \times 10^{-3}$ eV^2, $\theta_{12} = 34°$, $\theta_{13} = 0$, and $\theta_{23} = 45°$ and have set the Majorana phase equal to zero.

	M_{N_1} [MeV]	M_{N_2} [MeV]	m_ϕ [MeV]
N	1.2	1.2	0.85
I	1.4	1.4	1.0
N	100	1.2	0.85
I	100	1.3	0.97

TABLE II. The same notation as in Table I, but with complex ϕ.

	M_{N_1} [MeV]	M_{N_2} [MeV]	m_{ϕ_1} [MeV]	m_{ϕ_2} [MeV]
N	10^5	10^5	3.3	1
I	10^5	10^5	3.7	1
N	5.8	5.8	2.6	1.8
I	6.6	6.6	2.9	2.0

case of $m_{N_1} \neq m_{N_2}$, however, the cross section will depend on O. Moreover, as long as we neglect $\Delta m^2_{\rm sol}$, the cross section does not depend on the Majorana phase (neither for $m_{N_1} \neq m_{N_2}$ nor for $m_{N_1} = m_{N_2}$). Taking $\Lambda \simeq 200$ GeV and $\langle \sigma v_r \rangle \simeq 10^{-26}$ cm^3/s, we find from Eq. (20) that m_N should lie between 1 MeV and 1.5 MeV, depending on the ratio m^2_ϕ/m^2_N. In Tables I and II, examples of possible solutions are given. For simplicity, we have taken the arbitrary matrix O to have trivial format without mixing.

VI. CONCLUSION

In this paper, we propose a simple scenario, based on a single interaction term (Eq. (1)), where the dark matter candidate is an electroweak singlet scalar (named SLIM), which interacts with a Majorana fermion and a left-handed neutrino. This term generates left-handed neutrino masses through a one-loop diagram which can be directly related to either the SLIM pair annihilation cross section into neutrino pairs (or antineutrino pairs), as described in Eqs. (4) and (11) or, equivalently, to the SLIM relic density (see Eq. (6)). If the SLIM particle (our dark matter candidate) is a real scalar, the relation equaiton (6) leads to an upper bound of 10 MeV on the SLIM mass as well as a lower bound on its coupling (see Eq. (9)). Quite similarly, Eq. (11) constrains the mass splitting between the two components of the complex scalar field to lie in the MeV range.

Surprisingly enough, we obtain that in the real case, the favored range of the dark matter mass lies between a few MeV to 10 MeV, which corresponds to the light dark matter scenario. In principle, one could then explain the morphology of the 511 keV line emission in our galaxy [8] by extending Eq. (1) so as to include very weak-strength interactions with electrons.

A very exciting feature of the present scenario is that one can learn about the properties of dark matter by using precision measurements of the left-handed neutrino masses and the invisible decay modes of light mesons.

Finally the relations described in Eqs. (4) and (11) might indicate the existence of a low energy theory or new phenomena that are in principle testable in low energy experiments, but have not been detected in past experiments due to their lack of luminosity or sensitivity.

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