Monochromatic neutrinos generated by dark matter and the seesaw mechanism

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We study a minimal extension of the Standard Model where a scalar field is coupled to the right-handed neutrino responsible for the seesaw mechanism for neutrino masses. In the absence of other couplings, below 8 TeV the scalar A has a unique decay mode $A \to \nu\nu$, ν being the physical observed light neutrino state. Above 8 TeV (11 TeV), the 3-body (4-body) decay modes dominate. Imposing constraints on neutrino masses m_ν from atmospheric and solar experiments implies a long lifetime for A, much larger than the age of the Universe, making it a natural dark matter candidate. Its lifetime can be as large as 10^{29} seconds, and its signature below 8 TeV would be a clear monochromatic neutrino signal, which can be observed by ANTARES or IceCube. Under certain conditions, the scalar A may be viewed as a Goldstone mode of a complex scalar field whose vacuum expectation value generates the Majorana mass for ν_R. In this case, we expect the dark matter scalar to have a mass $\lesssim 10$ GeV.

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I. INTRODUCTION

The Standard Model (SM) of particle physics is now more than ever motivated by the recent discovery of the Higgs boson at both the ATLAS [1] and CMS [2] detectors. However, there are still two missing pieces in this elegant picture: the nature of dark matter (DM) and the origin of neutrino mass. Despite the fact that a window for low-mass dark matter candidates (below 100 TeV) seems favored by an upper bound coming for perturbative unitarity [3], no evidence has been found after many years of experimental searches [4]. On the other hand, the presence of new physics at an intermediate scale seems motivated by the stability of the Higgs potential [5,6], the seesaw mechanism [7,8], leptogenesis [9,10] or reheating processes. Added to the fact that a superheavy DM, or WIMPZILLA [11], could be produced by nonthermal processes and explain the DM density in the Universe, it seems natural to link the mechanism for generating a neutrino mass with the dark matter question in a coherent framework at an intermediate scale.

An intermediate scale (of order 10^{10} GeV) will never be reached by an accelerator on Earth. The 100 TeV collider is still a state-of-mind project, whereas the ILC can reach, at most, a beyond the SM (BSM) scale of 100 TeV through precision measurement. However, we know that energies as large as 10^{10} GeV are measured in ultrahigh-energy cosmic ray experiments like the Auger observatory [12]. Recently, the IceCube Collaboration claimed the detection of multi-PeV (10^6 GeV) events [13], giving the community some hope that an intermediate scale can be testable in the near future with these types of experiments.

In this paper we show that a high-energy neutrino signal can be associated with a long-lived scalar dark matter candidate. We show that this scalar can account for the dark matter in the Universe and moreover, its specific decay mode into two monochromatic neutrino states gives a clear signature detectable in present high-energy detectors like IceCube [13]. We then attempt to relate this candidate with the pseudo-Goldstone mode of a complex scalar field responsible for generating a Majorana mass in the right-handed sector through dynamical symmetry breaking at an intermediate scale.

IThe paper is organized as follows: After a description of the single scalar model we analyze in Sec. II, we compute its phenomenological consequences and detection prospects in Sec. III. In Sec. IV, we relate this scalar as the Goldstone mode associated with generating the right-handed neutrino mass, necessary for the seesaw mechanism. We draw our conclusions in Sec. V.

II. DARK MATTER AND A STANDARD SEESAW MECHANISM

A. The model

IAs a simple extension to the SM with a neutrino seesaw mechanism, we add a single real scalar field, A, coupled to the right-handed (sterile) sector. The Lagrangian can then be written as

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$$\mathcal{L} = \mathcal{L}_{\rm SM} + \mathcal{L}_\nu + \mathcal{L}_A \tag{1}$$

with

$$\mathcal{L}_\nu = -\left(\frac{1}{2}M^R + \frac{ih}{\sqrt{2}}A\right)\bar{\nu}_R^c\nu_R - \frac{y_{LR}}{\sqrt{2}}\bar{\nu}_L H\nu_R + \text{H.c.} \tag{2}$$

and

$$\mathcal{L}_A = -\frac{\mu_A^2}{2}A^2 - \frac{\lambda_A}{4}A^4 - \frac{\lambda_{HA}}{4}A^2H^2 + \frac{1}{2}\partial_\mu A\partial^\mu A, \tag{3}$$

where H represents the real part of the SM Higgs field. Here, we have simply assumed that the right-handed neutrino has a Majorana mass, M^R. We will explore a dynamical version of this extension in Sec. IV.

The scalar A is massive and couples to the SM Higgs, but does not itself get a vacuum expectation value (VEV). While there is no natural value for the mass scale M^R, demanding gauge coupling unification in different schemes of SO(10) breaking naturally leads to intermediate scales between 10^6 and 10^{14} GeV [14,15]. It seems then reasonable to expect that M^R will lie in this energy range if one embeds our model in a framework where one imposes unification of the gauge couplings. However, we will attempt to stay as general as possible.[1] In the context of very light scalar A, of the order of a keV (though not considered in the present work), some authors have looked at the effect of a decaying A on the CMB [17] and more recently the subleading effect of decays to photons [18].

B. The seesaw mechanism

Once symmetry breaking is realized, the mass states in the neutrino sector are mixed in the current eigenstate basis. Diagonalization of the mass matrix leads to the well-known seesaw mechanism. We can write the mass term

$$\mathcal{L}_\nu = -\frac{1}{2}\bar{n}\mathcal{M}n, \quad \text{with} \quad n = \begin{pmatrix} \nu_L + \nu_L^c \\ \nu_R + \nu_R^c \end{pmatrix} = \begin{pmatrix} n_1 \\ n_2 \end{pmatrix}$$

and

$$\mathcal{M} = \begin{pmatrix} 0 & m_D \\ m_D & M^R \end{pmatrix}, \tag{4}$$

with $m_D = y_{LR}v_H/\sqrt{2}$ ($v_H = 246$ GeV being the Higgs VEV). $\mathcal{M}$, being a complex symmetric matrix, can be diagonalized with the help of *one unitary matrix* U, $\mathcal{M} = UmU^T$, with

$$m = \begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix}. \tag{5}$$

From the diagonalization of $\mathcal{M}$,

$$m_1 = \frac{1}{2}\left[M^R - \sqrt{(M^R)^2 + 4m_D^2}\right] \simeq -\frac{m_D^2}{M^R} \simeq -\frac{y_{LR}^2 v_H^2}{2M^R},$$
$$m_2 = \frac{1}{2}\left[M^R + \sqrt{(M^R)^2 + 4m_D^2}\right] \simeq M^R, \tag{6}$$

and the eigenvectors N_1 and N_2

$$\begin{pmatrix} N_1 \\ N_2 \end{pmatrix} \simeq \begin{pmatrix} n_1 - \theta & n_2 \\ n_2 + \theta & n_1 \end{pmatrix} = \begin{pmatrix} \nu_L + \nu_L^c - \theta & (\nu_R + \nu_R^c) \\ \nu_R + \nu_R^c + \theta & (\nu_L + \nu_L^c) \end{pmatrix}, \tag{7}$$

where[2] $\tan 2\theta = -\frac{2m_D}{M^R}$, implying $\theta \simeq \sin\theta \simeq -\frac{m_D}{M^R} = -\frac{y_{LR}v_H}{\sqrt{2}M^R}$. Once the Lagrangian is expressed in terms of the physical mass eigenstates, one can compute the couplings generated by the symmetry breaking and their phenomenological consequences. m_1 corresponds to the mass of the Standard Model neutrino. We will consider $m_1 \lesssim 1$ eV from cosmological constraints through the rest of the paper.[3]

III. PHENOMENOLOGY

A. Generalities

To study the consequences of the model, we first rewrite the Lagrangian (2) in terms of the mass eigenstates, $N_{1,2}$:

$$\begin{aligned}\mathcal{L}_\nu = &-\frac{y_{LR}}{2\sqrt{2}}H(\bar{N}_1N_2 + \bar{N}_2N_1) - \frac{y_{LR}\theta}{\sqrt{2}}H(\bar{N}_2N_2 - \bar{N}_1N_1) \\ &-\frac{m_1}{2}\bar{N}_1N_1 - \frac{m_2}{2}\bar{N}_2N_2 \\ &-i\frac{h}{\sqrt{2}}A(\bar{N}_2\gamma^5N_2 - \theta\bar{N}_1\gamma^5N_2 - \theta\bar{N}_2\gamma^5N_1 + \theta^2\bar{N}_1\gamma^5N_1).\end{aligned} \tag{8}$$

A look at the Lagrangian (8) implies some obvious phenomenological consequences of the coupling of the scalar to the neutrino sector. First of all, the field N_2 is not stable through its decay $N_2 \to HN_1$ and cannot be the dark matter candidate as in the standard seesaw mechanism. Second, the scalar A is not stable, and its dominant decay mode for $M_A \lesssim 8$ TeV is $A \to N_1N_1$, as $M_{N_2} = m_2$ is of the order of M^R and is for now assumed to be heavier than A. When we include A as part of a dynamical mechanism for generating the mass M^R, we will see that the mass of A

[1] We note that a similar framework has been used in Ref. [16] to stabilize the Higgs potential up to GUT scale.

[2] Notice that N_1 and N_2 are Majorana-like particles.

[3] We neglect the flavor structure of the SM neutrino sector, as it does not affect our main conclusions.

may be highly suppressed relative to M^R, justifying *a posteriori* our assumption that $M_A < M_{N_1}$. Moreover, because the coupling of A to N_1 is of the order of $h\theta^2$, A is naturally long lived, and can be a good dark matter candidate, as we will see below. Its decay signature should be two ultra-energetic monochromatic neutrinos, which is a clear observable, and could be accessible to the present neutrino experiments like IceCube, ANTARES or SuperK, as we will see below. Above 8 TeV, the 3-body decay mode dominates, and above 11 TeV, the 4-body decay mode dominates. In this case, the signature of decay is then no longer a monochromatic signal but a neutrino spectrum, as we will see in the next section.

B. Neutrino flux

Before computing the flux of neutrinos expected on Earth from the decay of the scalar A, let us first check if it can be a reliable dark matter candidate, fulfilling $\tau_A > \tau_{\mathrm{Universe}} = 10^{17}$seconds.[4] The 2-body decay width for $A \to N_1 N_1$ is given by

$$\Gamma_A^2 = \frac{10^{-38} h^2}{8\pi}\left(\frac{m_1}{1\ \mathrm{eV}}\right)^2 \left(\frac{10^{10}\ \mathrm{GeV}}{M^R}\right)^2 M_A,$$

implying

$$\tau_A \sim 1.6 \times 10^{12} h^{-2} \left(\frac{1\ \mathrm{eV}}{m_1}\right)^2 \left(\frac{M^R}{10^{10}\ \mathrm{GeV}}\right)^2 \left(\frac{1\ \mathrm{TeV}}{M_A}\right)\ [\mathrm{s}], \tag{9}$$

where we have taken for reference $m_1 \lesssim 1$ eV as implied from the solar and atmospheric constraints on neutrino masses. As one can see, for a scalar mass of order 1 TeV, one can obtain lifetimes in excess of the age of the Universe for $M^R \gtrsim 10^{13} h$ GeV, making the scalar a potentially interesting dark matter candidate.

However, it is important to check multibody processes when $M_A > v_H$. Indeed,[5] the 3-body process $A \to N_1 N_1 H$ or the 4-body decay $A \to N_1 N_1 H H$, through the exchange of a virtual N_2 becomes dominant. Under the approximation of massless final states (largely valid when $M_A \gg m_h$), one obtains for the 3- and 4-body decay widths

$$\Gamma_A^3 = \frac{10^{-38} h^2}{3 \times 2^8 \pi^3}\left(\frac{m_1}{1\ \mathrm{eV}}\right)^2 \left(\frac{10^{10}\ \mathrm{GeV}}{M^R}\right)^2 \left(\frac{M_A}{v_H}\right)^2 M_A \tag{10}$$

and

$$\Gamma_A^4 = \frac{10^{-38} h^2}{9 \times 2^{14} \pi^5}\left(\frac{m_1}{1\ \mathrm{eV}}\right)^2 \left(\frac{10^{10}\ \mathrm{GeV}}{M^R}\right)^2 \left(\frac{M_A}{v_H}\right)^4 M_A, \tag{11}$$

which gives for the total width

$$\Gamma_A = \frac{10^{-38} h^2}{8\pi}\left(\frac{m_1}{1\ \mathrm{eV}}\right)^2 \left(\frac{10^{10}\ \mathrm{GeV}}{M^R}\right)^2 M_A \times \left[1 + \frac{1}{96\pi^2}\left(\frac{M_A}{v_H}\right)^2 + \frac{1}{18432\pi^4}\left(\frac{M_A}{v_H}\right)^4\right]. \tag{12}$$

From the expression above, we can easily compute that for $M_A \gtrsim 10\pi v_H \simeq 8$ TeV the 3-body, and then for $M_A \gtrsim 11$ TeV the 4-body, dominates the decay process. It will be interesting then to see what kind of constraints IceCube or ANTARES can impose on the parameter space of the model.[6]

The IceCube Collaboration recently released their latest analysis [13] concerning the (non)observation of ultrahigh-energetic neutrinos above 3 PeV. The PeV event rate expected at a neutrino telescope of fiducial volume $\eta_E V$ and nucleon number density n_N from a decaying particle of mass M_{DM}, mass density and width Γ_{DM} is [21]

$$\Gamma_{\mathrm{events}} = \eta_E V \times n_N \times \sigma_N \times L_{\mathrm{MW}} \times \frac{\rho_{\mathrm{DM}}}{M_{\mathrm{DM}}} \times \Gamma_{\mathrm{DM}} \simeq 3 \times 10^{59} \eta_E \frac{\Gamma_{\mathrm{DM}}}{M_{\mathrm{DM}}}\ \mathrm{years}^{-1}, \tag{13}$$

where σ_N ($\simeq 9 \times 10^{-34}$ cm^2 at 1 PeV) is the neutrino-nucleon scattering cross section, n_N is the nucleon number density in the ice ($n_N \simeq n_{\mathrm{Ice}} \simeq 5 \times 10^{23}/\mathrm{cm}^3$), L_{MW} is the rough linear dimension of our Galaxy (10 kpc) and $\rho_{\mathrm{DM}} \simeq 0.39\ \mathrm{GeV\,cm^{-3}}$ is the Milky Way dark matter density (taken near the Earth for the purpose of our estimate). The volume V is set to be 1 km^3, which is roughly the size of the IceCube detector, whereas the efficiency coefficient η_E depends on the energy of the incoming neutrino and lies in the range $10^{-2} - 0.4$ [22].

The neutrino-nucleon cross section is, however, highly dependent on the scattering energy, and the authors of Ref. [23] obtained

$$\sigma_N = 8 \times 10^{-36}\ \mathrm{cm}^2 \left(\frac{E_\nu}{1\ \mathrm{GeV}}\right)^{0.363} = 6 \times 10^{-36}\ \mathrm{cm}^3 \left(\frac{M_{\mathrm{DM}}}{1\ \mathrm{GeV}}\right)^{0.363}. \tag{14}$$

Implementing Eq. (14) in the expression (13), and adding an astrophysical factor $f_{\mathrm{astro}} \simeq 1$ corresponding

[4]A recent study [19] showed that taking into account the recent BICEP2 results, the real lifetime to consider should be $\gtrsim 10^{18}$ s. However, the constraints from IceCube are much stronger ($\tau_A \gtrsim 10^{28}$ seconds for M_A at the PeV scale), as we will see below.

[5]The authors want to thank the referee for having pointed out this possibility.

[6]An earlier analysis in another context was proposed in Ref. [20].

to the uncertainty in the distribution of dark matter in the galactic halo, one can write

$$\Gamma_{\rm events} = 1.5 \times 10^{57} \eta_E f_{\rm astro} \frac{\Gamma_{\rm DM}}{M_{\rm DM}^{0.637}} \ {\rm years}^{-1}, \tag{15}$$

where $\Gamma_{\rm DM}$ and $M_{\rm DM}$ are expressed in GeV. Noticing that there is no background from cosmological neutrinos at energies above 100 TeV, one can deduce the limit set by IceCube from the nonobservation of events above 3 PeV. IceCube took data during three years, so asking $3 \times \Gamma_{\rm events} \lesssim 1$, one obtains for $f_{\rm astro} = 1$

$$h^2 \left(\frac{M_A}{1\ {\rm GeV}}\right)^{4.363} \eta_E \lesssim 3.7 \times 10^{-3} \left(\frac{1\ {\rm eV}}{m_1}\right)^2 \left(\frac{M^R}{10^{10}\ {\rm GeV}}\right)^2. \tag{16}$$

If we take $M_A = 1$ PeV, one obtains $h \lesssim 8 \times 10^{-11}$ for $M^R \sim 10^{14}$ GeV, $\eta_E = 0.4$ and $m_1 = 1$ eV.

One can generalize our study to lower masses, down to the GeV scale, taking into account the combined constraints [24–26] from SuperK [27], ANTARES [28] and IceCube [29]. The limit on the lifetime of A as a function of M_A is depicted in Fig. 1. The resulting constraint in the (M_A, h) parameter plane is shown in Fig. 2 for different values of M_R. We see that natural values of M^R $(\gtrsim 10^{12})$ GeV lead to an upper limit on $h \lesssim 10^{-5}$, for $M_A > 1$ TeV.

We note that despite the fact that a dark matter source for the PeV events of IceCube is less motivated since the discovery of the third event "big bird," one can also compute the relation between h and M^R to observe the rate of 1 event per year for a 1 PeV dark matter candidate. We obtain from Eq. (15) $h \simeq 1.3 \times 10^{-10}$ for $M^R = 10^{14}$ GeV.

We also made a more detailed analysis, taking into account a simulated NFW galactic profile $\rho_{\rm NFW}$ for the Milky Way. Our result differs from the constraint with

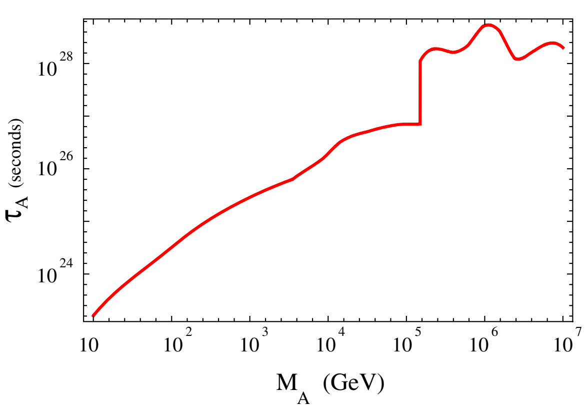

FIG. 1 (color online). Limit on the lifetime of A in seconds as a function of its mass M_A extracted from the combined constraints [24–26] from SuperK [27], ANTARES [28] and IceCube [29].

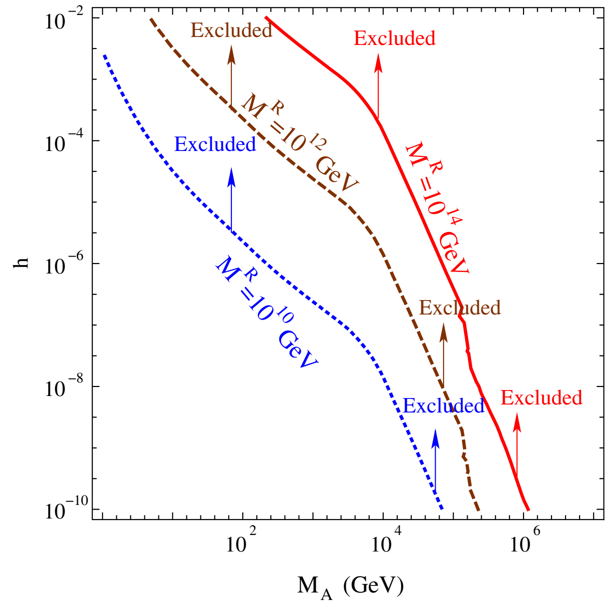

FIG. 2 (color online). Parameter space allowed in the plane (M_A, h), taking into account a combined analysis of IceCube and SuperK for different values of M^R and $m_1 = 1$ eV. The regions above the lines are excluded.

$f_{\rm astro} = 1$ only by a factor of a few (2–3). Indeed, compared to an annihilating scenario, due to the lack of quadratic enhancement in the signal, the role of the (better determined) local density is more prominent. We can thus anticipate little dependence of our conclusions on the specific galactic halo used for the analysis.

IV. DARK MATTER AND A DYNAMICAL SEESAW

A. The model

We would now like to ask whether or not the scalar A can be incorporated into a dynamical mechanism for generating neutrino seesaw masses. Instead of the coupling of A to ν_R in Eq. (2), let us couple the right-handed (sterile) sector to a complex scalar field $\Phi = S e^{ia/v_S}$. We assume that Φ is responsible for the breaking of some global symmetry so that S acquires a VEV. Here, we would like to stay as general as possible, and show that our framework can in fact be an illustration of any extension to the SM with dynamical breaking occurring at an intermediate scale. We now rewrite the Lagrangian as

$$\mathcal{L} = \mathcal{L}_{\rm SM} + \mathcal{L}_\nu + \mathcal{L}_\Phi \tag{17}$$

with

$$\mathcal{L}_\nu = -\frac{h}{\sqrt{2}} \bar{\nu}_R^c \Phi \nu_R - \frac{y_{LR}}{\sqrt{2}} \bar{\nu}_L H \nu_R + {\rm H.c.} \tag{18}$$

and

$$\mathcal{L}_\Phi = \frac{\mu_\phi^2}{2}|\Phi|^2 - \frac{\lambda_\Phi}{4}|\Phi|^4 - \frac{\lambda_{H\Phi}}{4}|\Phi|^2|H|^2 + \frac{1}{2}\partial_\mu\Phi\partial^\mu\Phi^*. \tag{19}$$

The Lagrangian above has a global $U(1)$ symmetry, under which the charges of ν_R, ν_L and Φ are 1, 1 and -2, respectively. After symmetry breaking generated by the fields H and Φ, one obtains in the heavy sector

$$\langle S\rangle = v_S = \frac{\mu_\Phi}{\sqrt{\lambda_\Phi}}; \qquad S = v_S + s; \qquad M_S = \sqrt{2}\mu_\Phi, \tag{20}$$

and we denote by A the argument of Φ after its magnitude is shifted by v_S. Note that our scalar field has been promoted to a Goldstone mode and is massless at tree level. The right-handed mass, M^R, is now given by $hv_S/\sqrt{2}$.

B. Breaking to a discrete symmetry

If our $U(1)$ symmetry were exact (prior to Φ picking up a VEV), M_A would remain massless to all orders in perturbation theory. In what follows, we will assume that the $U(1)$ symmetry is broken by nonperturbative effects down to a discrete Z_N symmetry. It is actually standard in string theory that all symmetries are gauged symmetries in the UV. Some of them are nonlinearly realized, in the sense that under gauge transformations one axion $\tilde\theta$ has a nonlinear transformation

$$A_\mu \to A_\mu + \partial_\mu\alpha, \qquad \Phi_i \to e^{iq_i\alpha}\Phi_i, \qquad \tilde\theta \to \tilde\theta + \alpha, \tag{21}$$

and the Lagrangian contains the Stueckelberg combination of a massive vector boson,

$$\frac{M^2}{2}(A_\mu - \partial_\mu\tilde\theta)^2. \tag{22}$$

Nonperturbative effects can generate operators of the form [30]

$$\frac{c_n}{M_P^{n-4}} e^{-2\pi N(t+i\frac{\tilde\theta}{2\pi})}\prod_i \Phi_i, \tag{23}$$

where t is a field which gets a VEV and where $S_{\rm inst} = 2\pi N(t + i\frac{\tilde\theta}{2\pi})$ can be interpreted as an instanton action. Nonperturbatively generated operators (23) are gauge invariant, provided that $\sum_i q_i = N$. The gauge field A_μ which eats the axion $\tilde\theta$ and the field t can be very heavy and decouple at low energy. At low energies, therefore, one gets an effective operator

$$e^{-\langle S_{\rm inst}\rangle}\frac{c_n}{M_P^{n-4}}\prod_{i=1}^n \Phi_i, \tag{24}$$

with c a numerical coefficient. At low energy, even though the $U(1)$ gauge symmetry is broken, one obtains a remnant Z_N symmetry. Due to its gauge origin, it satisfies anomaly cancellation conditions [31]. If the original gauge symmetry was anomaly free (which is realized if the axionic coupling to gauge fields as $\tilde\theta F\tilde F$ is absent), then anomalies have to be canceled. In particular, the mixed anomalies $Z_N SU(3)_c^2$, $Z_N SU(2)_L^2$ and $Z_N U(1)_Y^2$ anomalies have to vanish modulo N. For three generations of neutrinos, a simple candidate anomaly-free symmetry is Z_3. Then the lowest-order nonperturbative operator breaking $U(1)$ is

$$e^{-12\pi\langle t\rangle} c M_P(\Phi^3 + \bar\Phi^3) = e^{-12\pi\langle t\rangle} c M_P(2S^3 - 6SA^2). \tag{25}$$

This will generate a nonperturbative mass for the field A,

$$M_A^2 = 12 c v_S M_P e^{-12\pi\langle t\rangle}. \tag{26}$$

For moderate values of $\langle t\rangle$ this generates a large hierarchy for

$$\frac{M_A}{M_S} \sim e^{-6\pi\langle t\rangle}\sqrt{\frac{M_P}{v_S}}. \tag{27}$$

C. The signal

After the symmetry breaking, the Lagrangian (18) becomes

$$\begin{aligned}\mathcal{L}_\nu = &-\frac{h}{\sqrt{2}}s\bar N_2 N_2 + \frac{h\theta}{\sqrt{2}}s(\bar N_1 N_2 + \bar N_2 N_1) - \frac{y_{LR}}{2\sqrt{2}}H(\bar N_1 N_2 + \bar N_2 N_1) - \frac{y_{LR}\theta}{\sqrt{2}}H(\bar N_2 N_2 - \bar N_1 N_1)\\ &-\frac{m_1}{2}\bar N_1 N_1 - \frac{m_2}{2}\bar N_2 N_2 - i\frac{h}{\sqrt{2}}A(\bar N_2\gamma^5 N_2 - \theta\bar N_1\gamma^5 N_2 - \theta\bar N_2\gamma^5 N_1 + \theta^2\bar N_1\gamma^5 N_1).\end{aligned} \tag{28}$$

For $M_A \lesssim 8$ TeV, the dominant decay mode is the 2-body process $A \to N_1 N_1$, and the width of the A boson is now given by

$$\Gamma_A = \frac{10^{-38}}{4\pi}\left(\frac{m_1}{1\ \text{eV}}\right)^2\left(\frac{10^{10}\ \text{GeV}}{v_S}\right)^2 M_A,$$

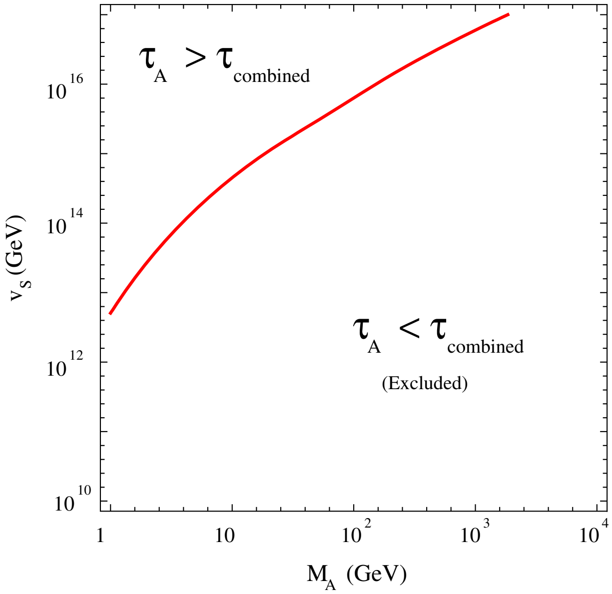

FIG. 3 (color online). Parameter space allowed by SuperK and Icecube, in the plane (M_A, v_S) for $m_1 = 1$ eV.

implying

$$\tau_A \sim 8 \times 10^{14} \left(\frac{1\ \mathrm{eV}}{m_1}\right)^2 \left(\frac{v_S}{10^{10}\ \mathrm{GeV}}\right)^2 \left(\frac{1\ \mathrm{GeV}}{M_A}\right)\ [\mathrm{s}]. \tag{29}$$

Since M^R is now proportional to hv_S, the decay rate becomes independent of h if we express it in terms of v_S and we are forced to consider sub-PeV masses for our dark matter candidate. As one can see, for relatively light Goldstone masses of order 1 GeV, one can obtain lifetimes in excess of the age of the Universe for $v_S > 10^{11}$ GeV, making this Goldstone mode an interesting dark matter candidate.

We show in Fig. 3 the parameter space allowed by the (non)observation of signals in neutrino telescopes. As surprising as it seems, we obtain quite reasonable values for v_S, compatible with GUT-like constructions.

V. CONCLUSION

In this work, we have shown that a dynamical model to generate Majorana neutrino masses naturally leads to the presence of a heavy quasistable pseudoscalar particle that can fill the dark matter component of the Universe and whose main decay mode into two ultra-energetic neutrinos is a clear signature observable by the IceCube detector. Our work is very general and can be embedded in many grand unified scenarios where the breaking of hidden symmetries appears at an intermediate scale.

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