Cold Dark Matter from heavy right-handed neutrino mixing

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We show that, within the seesaw mechanism, an almost decoupled right-handed (RH) neutrino species N_{DM} with mass $M_{DM} \gtrsim 100 \,\text{GeV}$ can play the role of Dark Matter (DM). The N_{DM} 's can be produced from nonadiabatic conversions of thermalized (source) RH neutrinos with mass M_S lower than M_{DM} . This is possible if a non-renormalizable operator is added to the minimal type I seesaw Lagrangian. The observed DM abundance can be reproduced for $M_{DM} \delta^{1/4} \sim 10^{-13} \text{ A}_{\text{eff}} \xi$, where Λ_{eff} is a very high energy new physics scale, $\delta \equiv (M_{DM} - M_S)/M_{DM}$ and $\xi \lesssim 1$ is a parameter determined by the RH neutrino ouplings.

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I. INTRODUCTION

The results from neutrino os
illation experiments represent a success for the seesaw mechanism $[1]$ $[1]$, the simplest way to understand why neutrinos are massive, yet so light compared to all other massive particles in the Standard Model (SM).

Indeed, within the seesaw, the atmospheric and the solar neutrino mass s
ales point to a high energy s
ale $\sim 10^{15}$ GeV compatible with grand unification and at the same time one can understand the observed large mixing angles. Moreover, neutrino os
illations support leptogen- $\text{e}\sin[2]$, an attractive way to explain the observed baryon asymmetry of the Universe and a dire
t onsequen
e of the seesaw me
hanism.

Despite the great progress made in recent years in deriving, especially from leptogenesis [3], interesting constraints on those seesaw parameters that escape the low energy experiments investigation, we still lack a way to probe the seesaw me
hanism. The main obsta
le is that, for natural hoi
es of the seesaw parameters, the heavy $right-handed(RH)$ neutrinos, predicted by the seesaw, are not expected to be detected at colliders, because they would be either too heavy or too weakly coupled. Moreover they usually de
ay very fast disappearing from the osmologi
al lore. If leptogenesis is the right explanation of the observed matter-antimatter asymmetry of the Universe, produ
ed from the CP violating de
ays of the RH neutrinos, this would be the only relic trace left over at present.

However, in this paper, we show that a weakly coupled RH neutrino species can play the role of cold DM. The scenario we present differs significantly from the one proposed in [\[5](#page-4-3)], where the lightest RH neutrino with a $\mathcal{O}(K$ eV) mass plays the role of warm DM, and neutrino Yukawa couplings are much smaller compared to charged leptons and quarks Yukawa ouplings. In our model, we assume that all RH neutrinos are heavy, with the lightest RH neutrino mass not lower than the electroweak scale. In this way, the neutrino Yukawa ouplings an be of the same order as for the other massive fermions.

II. FAILURE OF THE MINIMAL PICTURE

The (type I) seesaw mechanism $|1|$ is a minimal way to explain neutrino masses. The SM Lagrangian is extended adding a Yukawa interaction term between three RH neutrinos ν_R and the three left-handed doublets l via a Higgs doublet ϕ and a Majorana mass term M ,

$$
-\mathcal{L}_{Y+M} = \bar{l}_L \phi h \nu_R - \frac{1}{2} \overline{\nu_R^c} M \nu_R + h.c. , \qquad (1)
$$

where h is the matrix of the neutrino Yukawa couplings.

After ele
troweak symmetry breaking, indu
ed by the Higgs VEV v , the Yukawa interaction generates a Dirac mass term $m_D = hv$. In the seesaw limit, $M \gg m_D$, the spe
trum of mass eigenstates splits into three light neutrinos ν_i with masses given by the seesaw formula,

diag
$$
(m_1, m_2, m_3)
$$
 = $-U^{\dagger} m_D \frac{1}{M} m_D^T U^*$, (2)

where U is the leptonic mixing matrix, and into three heavy neutrinos N_i with masses $M_1 \leq M_2 \leq M_3$. These oin
ide, with very good approximation, with the eigenvalues of the Majorana mass matrix.

Neutrino os
illations experiments measure two neutrino mass-squared differences. For normal schemes one has $m_3^2 - m_2^2 = \Delta m_{\text{atm}}^2$ and $m_2^2 - m_1^2 = \Delta m_{\text{sol}}^2$, whereas
for inverted schemes one has $m_3^2 - m_2^2 = \Delta m_{\text{sol}}^2$ and m_2^2 $m_1^2 = \Delta m_{\text{atm}}^2$ $m_1^2 = \Delta m_{\text{atm}}^2$. For $m_1 \gg m_{\text{atm}} \equiv \sqrt{\Delta m_{\text{atm}}^2 + \Delta m_{\text{sol}}^2} =$ (0.050 ± 0.001) eV [6] the spectrum is quasidegenerate, while for $m_1 \ll m_{\text{sol}} \equiv \sqrt{\Delta m_{\text{sol}}^2} = (0.00875 \pm 0.00012) \text{ eV}$ [6] it is fully hierarchical (normal or inverted). For definiteness we will refer to the ase of normal s
hemes but all the dis
ussion applies to inverted s
hemes as well.

The RH neutrino de
ays an be onveniently des
ribed in terms of the decay parameters $K_i \equiv \Gamma_{Di}/H(T = M_i),$ where Γ_{Di} are the decay widths. These can be related to the neutrino masses introducing the effective neutrino masses, defined as $\widetilde{m}_i \equiv (m_D^{\dagger} m_D)_{ii}/M_i$, such that $K_i = \widetilde{\mathfrak{m}}_i/m_i$ assuming M to be $\widetilde{m}_i/m_{\star}$, where $m_{\star} \simeq 1.08 \times 10^{-3} \text{ eV}$. Assuming N_1 to be heavier than the Higgs boson,

from the LEP bound [7] one has $M_1 \gtrsim 115 \,\text{GeV}$ and the N_i lifetimes are then given by

$$
\tau_i = \frac{8\,\pi\,v^2}{\widetilde{m}_i\,M_i^2} \simeq \frac{5}{K_i} \left(\frac{\text{TeV}}{M_i}\right)^2 \times 10^{-13}\,\text{sec} \,. \tag{3}
$$

Let us now impose that one among the three RH neutrinos species N_i , plays the role of DM particle which we indicate with N_{DM} . This implies $\tau_{DM} \ge t_0 \simeq 4 \times 10^{17}$ sec, where t_0 is the age of the Universe. However, since the N_{DM} -decays would produce ordinary neutrinos, a much more stringent lower bound omes from neutrino telescopes $[8]$,

$$
\frac{\tau_{DM}}{t_0} \gtrsim \alpha \gg 1. \tag{4}
$$

In the range $M_{DM} \sim 10^{5\div 9}$ GeV, the AMANDA limits on neutrino flux implies $\alpha \sim 10^9$ [\[9](#page-4-7), 10], while in the range $M_{DM} \sim 10^{2 \div 5}$ GeV, where the atmospheric neutrino flux is observed, the lower bound is more relaxed. In any case, sin
e strong future improvements are expe
ted from the ICE-CUBE experiment, we will leave indicated the dependence on α in the following discussion [23]. From the relation [\(3\)](#page-1-0), this translates into an upper bound on the decay parameter K_{DM} (or equivalently on the effective neutrino mass \widetilde{m}_{DM}) given by

$$
K_{DM}(\widetilde{m}_{DM}/\text{eV}) \lesssim \frac{10^{-30(33)}}{\alpha} \left(\frac{\text{TeV}}{\text{M}_\text{i}}\right)^2. \tag{5}
$$

Moreover, imposing that the N_{DM} -abundance explains the measured DM ontribution to the energy density of the Universe, one finds a condition on $r_{DM} \equiv$ $(N_{N_{DM}}/N_{\gamma})_{\text{prod}}$, the ratio of the number of N_{DM} to the photon number at the time of the N_{DM} -production, occurring at temperatures higher than the electroweak phase transition,

$$
r_{DM} \sim 10^{-9} \left(\Omega_{\text{DM}} h^2\right) \frac{\text{TeV}}{M_{DM}} \sim 10^{-10} \frac{\text{TeV}}{M_{DM}}. \tag{6}
$$

Assuming that the correct value of r_{DM} is produced by some external mechanism, for example from inflaton deays, a trivial DM model is obtained if the ondition Eq. (5) is satisfied. Within such a scenario one can indifferently identify either N_1 or N_2 N_2 or N_3 with N_{DM} . The orthogonal seesaw matrix Ω [12], is a useful tool to parametrize the Dirac mass matrix m_D , such that

$$
m_D = U D_m^{1/2} \, \Omega \, D_M^{1/2} \,, \tag{7}
$$

with $D_m \equiv \text{diag}(m_1, m_2, m_3)$ and $D_M \equiv \text{diag}(M_1, M_2, M_3)$. The effective neutrino masses The effective neutrino masses an then be expressed as linear ombinations of the neutrino masses $\widetilde{m}_i = \sum_h m_h |\Omega_{hi}|^2$ and one easily obtains $\widetilde{m}_i \geq m_1$. Therefore, the upper bound Eq. [\(5\)](#page-1-1) applies to m_1 as well, implying hierarchical light neutrinos. It also implies that Ω has to be close to the special form

$$
\begin{pmatrix}\n1 & 0 & 0 \\
0 & \cos \omega & \sin \omega \\
0 & -\sin \omega & \cos \omega\n\end{pmatrix}, \n(8)
$$

or to those other two obtained by column cyclic permutation. Therefore, assuming exactly one of these three forms for the orthogonal matrix, the condition Eq. (6) is fulfilled only assuming some mechanism for the N_{DM} produ
tion based on physi
s beyond the type I seesaw SM extension. Even allowing small deviations from these special forms, one undergoes a severe obstacle within the type I seesaw. Indeed one can think of different processes producing the N_{DM} -abundance, such as inverse decays or s
atterings involving the top quark or gauge bosons. However, in all cases one has approximately $r_{DM} \sim K_{DM}$ and it would then be impossible to satisfy simultaneously the two requirements Eq. [\(5\)](#page-1-1) and Eq. [\(6\)](#page-1-2).

Let us consider a particular example that clearly shows such a difficulty but that at the same time, as we will see, will suggest a solution relying on a simple and reasonable extension of the type I seesaw lagrangian.

We investigate the possibility that the N_{DM} production is induced by the mixing of N_{DM} with one of the other two RH neutrinos acting as a source, and that we indicate with N_S . Notice that N_S has necessarily a thermal abundan
e if the reheat temperature is approximately higher than M_S . This is because there cannot be more than one RH neutrino species with $\widetilde{m}_i \lesssim m_{\star}$.

For definiteness we can assume that N_{DM} and N_S are the two lightest RH neutrinos and hen
e there are only two possibilities: either $M_{DM} = M_1$ and $M_S = M_2$ or vice-versa. In this case N_3 does not play any role in the N_{DM} -production but it is necessary to reproduce correctly the neutrino masses.

This s
enario is realized hoosing the following form for the orthogonal matrix

$$
\Omega' = \begin{pmatrix}\n\sqrt{1 - \varepsilon^2} & -\varepsilon & 0 \\
\varepsilon & \sqrt{1 - \varepsilon^2} & 0 \\
0 & 0 & 1\n\end{pmatrix},
$$
\n(9)

representing a perturbation, with $\cos \omega = 1$, of the special form in Eq. (8) . Here the prime index indicates that we are reexpressing Ω into a basis where the RH neutrino mass term is still diagonal but in a way that M_{DM} is always the first eigenvalue and M_S the second eigenvalue. Notice that we can choose ε real and for convenience positive. Moreover notice that the choice $\cos \omega = 1$ is not restrictive. Indeed, in any case a value $\cos \omega \neq 1$ would not be relevant for the DM production but notice that it would be important if one simultaneously imposes successful leptogenesis from N_S N_S decays, a possibility that will be discussed elsewhere [13].

In order to describe the RH neutrino mixing, it is convenient to work in the "Yukawa basis", where the Yukawa interaction term is diagonal. This can be diagonalized by mean of a bi-unitary transformation, $D_h \equiv$ $diag(h_A, h_B, h_C) = V_L^{\dagger} h U_R$. The RH neutrino mixing matrix U_R can be found considering that it diagonalizes $h^{\dagger} h$, namely $U_R^{\dagger} (h^{\dagger} h) U_R = \text{diag}(h_A^2, h_B^2, h_C^2)$. Then, from the expression Eq. (7) , one can see that our choice for Ω' simply results into

$$
U_R = \begin{pmatrix} \cos \theta & -\sin \theta & 0 \\ \sin \theta & \cos \theta & 0 \\ 0 & 0 & 1 \end{pmatrix}, \qquad (10)
$$

with $\sin \theta \simeq \varepsilon \sqrt{M_S/M_{DM}}$ and into

$$
h_A \simeq \frac{\sqrt{m_1 M_{DM}}}{v}, h_B \simeq \frac{\sqrt{m_{sol} M_S}}{v}, h_C \simeq \frac{\sqrt{m_{atm} M_3}}{v}.
$$
\n(11)

This clearly shows that though N_3 does not mix, it is necessary to reproduce the atmospheric neutrino mass scale. Imposing the condition [\(5\)](#page-1-1), one can see that ε has to be tiny. Indeed one has

$$
\widetilde{m}_1 \simeq m_1 + m_{\rm sol} \, |\varepsilon|^2 \,, \tag{12}
$$

and therefore the upper bound Eq. [\(5\)](#page-1-1) translates into the upper bounds $[24]$

$$
\frac{m_1}{\rm eV} \lesssim \frac{10^{-33}}{\alpha} \, \left(\frac{\rm TeV}{\rm M_{DM}}\right)^2 \,, \quad |\varepsilon| \lesssim \frac{10^{-16}}{\sqrt{\alpha}} \, \left(\frac{\rm TeV}{M_{DM}}\right) \,. \tag{13}
$$

This implies a hierarchical light neutrino spectrum and a tiny mixing angle between the two lightest RH neutrinos. The description of the production of the N_{DM} abundan
e pro
eeds very similarly to the ase of light active-sterile neutrino oscillations [15] and in particular to the case described in [16], where transitions occur in the non-adiabatic regime as it will prove to be in our case. Let us write down the hamiltonian for the two lightest RH neutrinos in the Yukawa basis. This will be the sum of two terms: a pure kinetic term and a second term accounting for matter effects described by a potential that in the Yukawa basis is diagonal and given by $[17]$

$$
V_I \sim h_I^2 T^2 / (8 k) \qquad (I = A, B), \qquad (14)
$$

in the approximation of ultrarelativistic neutrinos, implying $E \sim k$ and $T \gg M_S/3$. Notice that in any case for $T \lesssim M_S$ the N_S-abundance is exponentially suppressed and the N_{DM} -production would stop anyway. In order to further simplify the problem, we also employ a mono
hromati approximation where all neutrinos have the same mean energy value $k \sim 3T$. As usual, we can subtract from the hamiltonian a term proportional to the identity, irrelevant in neutrino os
illations. Therefore, in the Yukawa basis, the relevant hamiltonian can be recast as

$$
\Delta H = \frac{\Delta M^2}{12 T} \begin{pmatrix} -\cos 2\theta + (v_A - v_B) & \sin 2\theta \\ \sin 2\theta & \cos 2\theta - (v_A - v_B) \end{pmatrix}
$$
(15)

where we defined $v_I \equiv T^2 h_I^2/(4 \Delta M^2)$ and $\Delta M^2 \equiv M_S^2$ M_{DM}^2 . Approximating $\cos 2\theta \simeq 1$, one can see that there is a resonan
e at a temperature

$$
T_{\rm res} \simeq 2 \sqrt{\frac{\Delta M^2}{h_A^2 - h_B^2}} \simeq 2 \frac{\sqrt{-\Delta M^2}}{h_B}, \qquad (16)
$$

,

only if $\Delta M^2 < 0$, i.e. only if $M_1 = M_S < M_{DM} = M_2$. Using the Eq. (11) , T_{res} can be conveniently recast as

$$
T_{\rm res} \simeq 10^7 \, M_{DM} \sqrt{\frac{v}{M_S} \left(1 - \frac{M_S^2}{M_{DM}^2}\right)}.
$$
 (17)

If $M_{DM} \gtrsim 2 M_S$ one has $T_{\rm res} \simeq 10^7 \, M_{DM} \, \sqrt{v/M_S}$. In this case, introducing $z_{\text{res}} \equiv M_{DM}/T_{\text{res}} \simeq 10^{-7} \sqrt{M_S/v}$, one can envisage a problem. The $N_S{\rm 's}$ thermalize for $z_{\text{eq}} \simeq (6/K_S)^{1/3} \simeq 0.8$ [\[4](#page-4-16)]. Imposing $z_{\text{res}} > z_{\text{eq}}$ leads to an unacceptably large values of M_S , M_D and of the reheat temperature. Therefore, unless one assumes an initial thermal abundan
e, one is for
ed to onsider the degenerate limit, for $\delta \equiv (M_{DM}-M_S)/M_{DM} \ll 1$. In this limit one now obtains $T_{\text{res}} \simeq 10^7 \, M_{DM} \, \delta^{1/2} \, \sqrt{v/M_{DM}}$ and $z_{\rm res} \simeq 10^{-7} \, \delta^{-1/2} \, \sqrt{M_{DM}/v}$. For $\delta \lesssim 10^{-13} \, M_{DM}/\text{TeV}$, this time one can have $z_{\text{res}} \gtrsim z_{\text{eq}}$. Therefore, the degenerate limit has to be onsidered as a more attra
tive option.

Be
ause of the tiny mixing angle the transitions at the resonance occur in the nonadiabatic regime. Indeed let us calculate the adiabaticity parameter at the resonance,

$$
\gamma_{\rm res} \equiv \frac{1}{2 \,\dot{\theta}_m \,\ell_m} \bigg|_{\rm res} = \sin^2 2\theta \, \frac{|\Delta M^2|}{6 \, T_{\rm res} \, H_{\rm res}} \,. \tag{18}
$$

Here $H_{\text{res}} \simeq 1.66 \sqrt{g_{\star}} T_{\text{res}}^2 / M_{\text{Pl}}$ is the value of the expansion rate at the resonan
e. Using the onditions Eq. [\(13\)](#page-2-1) and Eq. [\(4\)](#page-1-5), one obtains the upper bound $\gamma_{\rm res} \lesssim 10^{-26} \, \left({\rm TeV}/M_{DM} \right)^2$. The N_{DM} -abundance r_{DM} can then be calculated as the fraction of N_S 's that is converted into N_{DM} . This is approximately given by the Landau-Zener formula,

$$
r_{N_{DM}} \sim \frac{N_{DM}}{N_S} \sim (1 - e^{-\frac{\pi}{2}\gamma_{\text{res}}}) \simeq \frac{\pi}{2}\gamma_{\text{res}}.
$$
 (19)

Comparing with the ondition Eq. [\(6\)](#page-1-2), it is evident that neutrino mixing between heavy RH neutrinos annot produce the right N_{DM} abundance, at least not within a minimal type I seesaw extension of the SM. This conclusion is confirmed by more precise calculations beyond the Landau-Zener approximation.

III. A WAY-OUT FROM NONRENORMALIZABLE OPERATORS

Let us onsider the possibility that adding higher dimensional effective operators to the minimal type I seesaw Lagrangian Eq. (1) , while not affecting neutrino masses and mixing, enhances the N_{DM} -production from neutrino mixing. In particular let us consider the follow-ing dim-five effective operator [\[25](#page-4-17)]

$$
\mathcal{L}_{\text{eff}} \propto \frac{\lambda_{AB}}{\Lambda_{\text{eff}}} |\Phi|^2 \bar{N}_A^c N_B , \qquad (20)
$$

where Φ is the usual Higgs field, λ is a dimensionless coupling matrix and Λ_{eff} is an unspecified very high energy new physi
s s
ale that we treat as a free parameter.

This operator yields a new contribution to 'matter effects' into the hamiltonian [13], that in the Yukawa basis an be written as

$$
H_{\text{eff}} \simeq \frac{T^2}{12 \,\Lambda_{\text{eff}}} \,\lambda_{IJ} \,. \tag{21}
$$

This result follows from the omputation of the temperature dependent finite real part of the RH neutrino selfenergy $[17]$:

$$
\text{Re}[\Sigma_N(T)] = \frac{\lambda_{IJ}}{\Lambda} \int \frac{d^4q}{(2\pi)^3} \delta(q^2 - m_\Phi^2) n_b(q), \qquad (22)
$$

where $n_b(q) = \frac{1}{e^{|q \cdot u|} - 1}$ is the Bose-Einstein distribution with u being the four-velocity of the thermal bath. Assuming zero Higgs mass one then immediately dedu
e orresponding orre
tion to the Hamiltonian [\(21\)](#page-3-0).

We can reasonably assume that $h_B^2 \gg T_{\text{res}}/\Lambda_{\text{eff}}$. In this way in the Yukawa basis the total intera
tion term is approximately still diagonal and with the same eigenvalues. The relevant hamiltonian des
ribing neutrino os illations an then be written as

$$
\Delta H^{\text{eff}} \simeq \frac{\Delta M^2}{12 T} \begin{pmatrix} -v_B & \sin 2\theta + v_{\text{eff}}^{AB} \\ \sin 2\theta + v_{\text{eff}}^{AB} & v_B \end{pmatrix}, \quad (23)
$$

where we introduced $v_{\text{eff}}^{IJ} \equiv T^3 \lambda_{IJ} / (2 \Delta M^2 \Lambda_{\text{eff}})$. Notice that the resonance condition on the temperature, Eq. [\(16\)](#page-2-2), does not hange. However, now the mixing angle is different and receives a contribution from the off-diagonal terms in H_{eff} , such that $\sin 2\theta_{\text{eff}} \simeq v_{\text{eff}}^{AB}$.

Imposing again that mixing is responsible for the DM production, since we know that the mixing angle θ induced by the Yukawa coupling h_A is by far too small to play any role, it an be assumed to be exa
tly zero. This is a good feature sin
e otherwise one ould have ob je
ted that radiative orre
tions ould indu
e a large value anyway, spoiling the stability of N_{DM} on cosmological scales. However, if it is exactly zero, one can invoke some symmetry that protects it from radiative corrections.

Therefore, the adiabaticity parameter can now be written as

$$
\gamma_{\rm res}^{\rm eff} \simeq \sin^2 2\theta_{\rm eff} \, \frac{|\Delta M^2|}{6 \, T_{\rm res} \, H_{\rm res}} \simeq \frac{\sqrt{|\Delta M^2|} \, M_{\rm Pl}}{5 \, \Lambda_{\rm eff}^2 \, \xi^2} \,, \qquad (24)
$$

where we used the Eq. (16) for T_{res} and defined $\xi \equiv g_{\star}^{1/4} h_B^{3/2} / \lambda_{AB}$. Using again the Landau-Zener approximation for an estimation of the N_{DM} abundance, $r_{N_{DM}} \sim \gamma_{\rm res}$, and imposing again the condition Eq. [\(6\)](#page-1-2), we obtain the ondition

$$
M_{DM} \,\delta^{\frac{1}{4}} \sim 10^{-13} \,\Lambda_{\text{eff}} \,\xi \,. \tag{25}
$$

It is easy to verify that the assumption is h_I^2 \gg T_{res}/ Λ _{eff}, translates into a condition

 M_S \gg 10⁻² GeV $g_{\star}^{1/3}$ $\delta^{2/3}/\lambda_{AB}^{4/3}$, easily verified except for tiny values of λ_{AB} . Notice also that using the Eq. [\(11\)](#page-2-0), one can recast $\xi \sim (10^{-9}/\lambda_{AB}) (M_S/\text{TeV})^{3/4}$ From the condition Eq. (25) , one then finds in the hierarchical case, i.e. $M_{DM} \geq 2 M_S$,

$$
M_S \lesssim \left(\frac{\Lambda_{\text{eff}}}{10^{13} \,\text{TeV}}\right)^4 \left(\frac{10^{-9}}{\lambda_{AB}}\right)^4 \,\text{TeV},\qquad(26)
$$

showing that in order not to satisfy $M_S \gtrsim 100 \,\text{GeV}$ the ouplings annot be too large. On the other hand in the more interesting degenerate limit $(\delta \ll 1)$ one finds

$$
M_{DM} \gg \left(\frac{\Lambda_{\text{eff}}}{10^{13} \text{ TeV}}\right)^4 \left(\frac{10^{-9}}{\lambda_{AB}}\right)^4 \text{ TeV},\qquad(27)
$$

showing, onversely, that in order not to have too large values of M_{DM} the couplings cannot be too small. Notice that too large values $\log(M_{\rm DM}/\rm{TeV}) \lesssim 5 \div 8$ would spoil the cosmologically stability of N_{DM} , leading to unobserved neutrino fluxes at neutrino telescopes. Indeed in this ase the nonrenormalizable operator and the mixing with M_S would induce too fast decays of the N_{DM} 's into Higgs and leptons [26]. For $\Lambda_{\text{eff}} \sim M_{GUT} \div M_{Pl}$ one has then $\lambda_{AB} \gtrsim 10^{-13 \div -10}$. The smallness of λ_{AB} can be explained in two ways. In the case when $\Lambda_{\text{eff}} \sim M_{GUT}$ the operator [\(20\)](#page-2-3) an be generated radiatively from the coupling to the GUT scale particles. For example, one can assume the Yukawa coupling (with the strength h) between RH neutrino, Higgs and heavy $(m \sim M_{GUT})$ fermion. This oupling generates at one loop the operator [\(20\)](#page-2-3) after heavy fermion is integrated out. The values of λ_{AB} are, therefore, given by $h^2(T_{\text{res}})$ and, if $h(T_{\text{res}}) \gtrsim 10^{-4 \div -5}$, they come out naturally in the desired region. Alternatively, if the operator [\(20\)](#page-2-3) is generated gravitationally ($\Lambda_{\text{eff}} \sim M_{Pl}$) the smallness of the coefficients λ_{AB} can be explained in the models where the effective value of M_{Pl} in the early universe is dif-ferent from its present value (e.g. see Refs.[\[19](#page-4-19)]). However, the onsequent de
ay hannels at present should be estimated with $\lambda_{AB} \sim 1$. A detailed analysis of the constraints from decays will be presented elsewhere [13], however, it is remarkable that the me
hanism is viable for reasonable values of the involved parameters.

IV. CONCLUSIONS

We presented a new s
enario where the role of DM is played by heavy RH neutrinos. The s
enario is based on a me
hanism where the DM RH neutrinos are produ
ed through mixing enhan
ed by the additional presen
e of higher dimensional effective operators into the usual type I seesaw Lagrangian. The mechanism relies crucially on the fa
t that is ne
essary to onvert just a very small fra
tion of the sour
e RH neutrinos into the DM RH neutrinos. In this way the additional operator has the effect of enhancing the mixing without spoiling any other sucessful feature of the type I seesaw me
hanism and at

the same time preserving the DM RH neutrinos stability on osmologi
al times. A straightforward predi
tion of the me
hanism is that the lightest neutrino mass has to vanish. It also seems quite general that the DM RH neutrinos decay and this could lead to signatures in cosmic rays. The recent detected excess of positrons in the HEAT and PAMELA experiments have been interpreted as due to de
aying DM parti
les with a mass higher than [300](#page-4-20) GeV and a lifetime of approximately $\tau_{DM} \sim 10^{26}$ sec [20]. Therefore, our mechanism seems to have the right features to explain this excess. These results are quite interesting sin
e not only they are fully ompatible with our model but also be
ause the value for the life time orresponds to the saturation of the lower bound Eq. [4](#page-1-5) from the AMANDA data when $M_{DM} \sim 10^{5 \div 9}$ GeV and

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a signal should be expe
ted from the ICE CUBE experiment.

It should also be noticed that the special orthogonal form Eq. (9) predicted by the mechanism corresponds [\[21](#page-4-21)] to a particular sequential dominated model [22]. Therefore, the proposed s
enario for the solution of the DM onundrum restri
ts remarkably the seesaw parameter spa
e, providing a potential smoking gun for the seesaw me
hanism.

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- [23] Constraints from γ -rays from decaying Dark Matter in the Milky Way halo give similar or even stronger bound depending on M_{DM} [\[11](#page-4-24)].
- [24] In the exact limit $m_1 = 0$ the eigenvalue $h_A = 0$. Moreover, in this limit, plugging $\Omega' = I$ (cf. (Eq. [9\)](#page-1-6)) into the Eq. [\(7\)](#page-1-4), one can immediately see that the Yukawa coupling matrix, in the basis where the RH neutrino mass matrix is diagonal, contains a vanishing column corresponding to the DM RH neutrino. This is quite an obvious result, since it corresponds to impose $\tau_{DM} = 0$ that implies $h_{\alpha DM} = 0$. Therefore, the model corresponds to a special textured form for the Yukawa couplings equivalent to impose a particular symmetry such that the lagrangian is invariant under a proper transformation of the neutrino fields.
- [25] The idea that this operator could enhance RH neutrino oscillations was presented preliminarily in [\[14](#page-4-23)]. It has then been also onsidered for non-resonant produ
tion of the sterile DM neutrinos in $[18]$ $[18]$.
- [26] This can be also regarded saying that the nonrenormalizable operator breaks the symmetry that brings the Yukawa oupling matrix to the spe
ial textured form which is needed to have N_{DM} cosmologically stable.