Cold Dark Matter from heavy right-handed neutrino mixing

Alexey Anisimov
Deutsches Elektronen-Synchrotron DESY, 22603 Hamburg, Germany

Pasquale Di Bari
Dipartimento di Fisica Galileo Galilei Via Marzolo 8, I-35131 Padua, Italy
(Dated: August 24, 2009)

We show that, within the see-saw 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 non-adiabatic 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 see-saw lagrangian. The observed DM abundance can be reproduced for $M_{DM}\,\delta^{1/4} \sim 10^{-13}\,\Lambda_{\rm eff}\,\xi$, where $\Lambda_{\rm eff}$ 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 couplings.

PACS numbers: 14.60.St, 95.35.+d

I. INTRODUCTION

The results from neutrino oscillation experiments represent a success for the seesaw mechanism [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 scales point to a high energy scale $\sim 10^{15}\,\text{GeV}$ compatible with grand unification and at the same time one can understand the observed large mixing angles. Moreover, neutrino oscillations support leptogenesis [2], an attractive way to explain the observed baryon asymmetry of the Universe and a direct consequence of the seesaw mechanism.

Despite the great progress made in the last 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 mechanism. The main obstacle is that, for natural choices of the seesaw parameters, the heavy 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 decay very fast disappearing from the cosmological lore. If leptogenesis is the right explanation of the observed matter-antimatter asymmetry of the Universe, produced from the *CP* violating decays 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 that one proposed in [5], where the lightest RH neutrino with a $\mathcal{O}(\text{KeV})$ mass plays the role of warm DM, and neutrino Yukawa couplings are much smaller compared to charged leptons and quarks Yukawa couplings. 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 couplings can 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 electroweak symmetry breaking, induced by the Higgs VEV v, the Yukawa interaction generates a Dirac mass term $m_D = h\,v$. In the seesaw limit, $M \gg m_D$, the spectrum of mass eigenstates splits into three light neutrinos ν_i with masses given by the seesaw formula,

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

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 coincide, with very good approximation, with the eigenvalues of the Majorana mass matrix.

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

The RH neutrino decays can be conveniently described in terms of the decay parameters $K_i \equiv \widetilde{\Gamma}_{Di}/H(T = M_i)$, where $\widetilde{\Gamma}_{Di}$ 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 $K_i = \widetilde{m}_i/m_\star$, 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 $M_1 \gtrsim 115\,\mathrm{GeV}$ and the N_i life-times are then given by

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

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 $\tau_{DM} \geq t_0 \simeq 4\times 10^{17}\,\mathrm{sec}$, 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 comes from neutrino telescopes,

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

In the range $M_{DM} \sim 10^{5\div 9}\,\mathrm{GeV}$, the AMANDA limits on neutrino flux implies $\alpha \sim 10^9$ [8, 9], while in the range $M_{DM} \sim 10^{2\div 5}\,\mathrm{GeV}$, where the atmospheric neutrino flux is observed, the lower bound is more relaxed. In any case, since strong future improvements are expected from the ICE-CUBE experiment, we will leave indicated the dependence on α in the following discussion. From the relation (3), this translates into an upper bound on the decay parameter K_{DM} (or equivalently on the effective neutrino mass $\widetilde{m}_{DM}$) given by

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

Moreover, imposing that the N_{DM}-abundance explains the measured DM contribution to the energy density of the Universe, one finds a condition on $r_{DM} \equiv (N_{N_{DM}}/N_\gamma)_{\mathrm{prod}}$, 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 electro-weak phase transition,

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

Assuming that the correct value of r_{DM} is produced by some external mechanism, for example from inflaton decays, a trivial DM model is obtained if the condition Eq. (5) is satisfied. Within such a scenario one can indifferently identify either N_1 or N_2 or N_3 with N_{DM}. The orthogonal seesaw matrix Ω [10], is a useful tool to parameterize the Dirac mass matrix m_D, such that

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

with $D_m \equiv \mathrm{diag}(m_1, m_2, m_3)$ and $D_M \equiv \mathrm{diag}(M_1, M_2, M_3)$. The effective neutrino masses can then be expressed as linear combinations 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) applies to m_1 as well, implying hierarchical light neutrinos. It also implies that Ω has to be close to the special form

$$\begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos\omega & \sin\omega \\ 0 & -\sin\omega & \cos\omega \end{pmatrix}\,, \tag{8}$$

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}-production based on physics 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 scatterings involving the top quark or gauge bosons. However, in all cases one has approximately $r_{DM} \sim K_{DM}$ and it would be then impossible to satisfy simultaneously the two requirements Eq. (5) and Eq. (6).

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 abundance 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 hence 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 scenario is realized choosing the following form for the orthogonal matrix

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

representing a perturbation, with $\cos\omega = 1$, of the special form in Eq. (8). Here the prime index indicates that we are re-expressing 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 decays, a possibility that will be discussed elsewhere [11].

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, $D_h \equiv \mathrm{diag}(h_A, h_B, h_C) = V_L^\dagger\,h\,U_R$. The RH neutrino mixing matrix U_R can be found considering that it diagonalizes $h^\dagger h$, namely $U_R^\dagger\,(h^\dagger\,h)\,U_R = \mathrm{diag}(h_A^2,h_B^2,h_C^2)$. Then, from the expression Eq. (7), one can see that our choice for Ω' simply results into

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

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

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

This clearly shows that though N_3 does not mix, it is necessary to reproduce the atmospheric neutrino mass scale. Imposing the condition (5), 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) translates into the upper bounds [20]

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

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}-abundance proceeds very similarly to the case of light active-sterile neutrino oscillations [12] and in particular to the case described in [13], 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 [14]

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

in the approximation of ultra-relativistic 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 monochromatic approximation where all neutrinos have the same mean energy value $k\sim 3\,T$. As usual, we can subtract from the hamiltonian a term proportional to the identity, irrelevant in neutrino oscillations. Therefore, in the Yukawa basis, the relevant hamiltonian can be recast as

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

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 resonance at a temperature

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

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

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

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

Because of the tiny mixing angle the transitions at the resonance occur in the non-adiabatic regime. Indeed let us calculate the adiabaticity parameter at the resonance,

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

Here $H_{\rm res} \simeq 1.66\,\sqrt{g_\star}\,T_{\rm res}^2/M_{\rm Pl}$ is the value of the expansion rate at the resonance. Using the conditions Eq. (13) and Eq. (4), one obtains the upper bound $\gamma_{\rm res} \lesssim 10^{-26}\,(\mathrm{TeV}/M_{DM})^2$. 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,

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

Comparing with the condition Eq. (6), it is evident that neutrino mixing between heavy RH neutrinos cannot produce the right N_{DM}-abundance, not at least 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 NON-RENORMALIZABLE OPERATORS

Let us consider 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 following dim-five effective operator[21]

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

where Φ is the usual Higgs field, λ is a dimensionless coupling matrix and $\Lambda_{\rm eff}$ is an unspecified very high energy new physics scale that we treat as a free parameter.

This operator yields a new contribution to 'matter effects' into the hamiltonian [11], that in the Yukawa basis can be written as

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

This result follows from the computation of the temperature dependent finite real part of the RH neutrino self-energy [14]:

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

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 deduce corresponding correction to the Hamiltonian (21).

We can reasonably assume that $h_B^2 \gg T_{\rm res}/\Lambda_{\rm eff}$. In this way in the Yukawa basis the total interaction term is approximately still diagonal and with the same eigenvalues. The relevant hamiltonian describing neutrino oscillations can then be written as

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

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

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 can be assumed to be exactly zero. This is a good feature since otherwise one could have objected that radiative corrections could induce 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 be now written as

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

where we used the Eq. (16) for $T_{\rm res}$ 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), we obtain the condition

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

It is easy to verify that the assumption $h_B^2 \gg T_{\rm res}/\Lambda_{\rm eff}$, translates into a condition $M_S \gg 10^{-2}\,{\rm GeV}\,g_\star^{1/3}\,\delta^{2/3}/\lambda_{AB}^{4/3}$, easily verified except for tiny values of λ_{AB}. Notice also that using the Eq. (11), one can recast $\xi \sim (10^{-9}/\lambda_{AB})\,(M_S/{\rm TeV})^{3/4}$. From the condition Eq. (25), one then finds in the hierarchical case, i.e. $M_{DM} \gtrsim 2\,M_S$,

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

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

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

showing, conversely, 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 case the non-renormalizable operator and the mixing with M_S would induce too fast decays of the N_{DM}'s into Higgs and leptons[22]. For $\Lambda_{\rm eff} \sim M_{GUT} \div M_{Pl}$ 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 $\Lambda_{\rm eff} \sim M_{GUT}$ the operator (20) can 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 coupling generates at one loop the operator (20) after heavy fermion is integrated out. The values of λ_{AB} are, therefore, given by $h^2(T_{\rm res})$ and, if $h(T_{\rm res}) \gtrsim 10^{-4\div 5}$, they come out naturally in the desired region. Alternatively, if the operator (20) is generated gravitationally ($\Lambda_{\rm eff} \sim M_{Pl}$) 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 different from its present value (e.g. see Refs.[19]). However, the consequent decay channels at present should be estimated with $\lambda_{AB} \sim 1$.** A detailed analysis of the constraints from decays will be presented elsewhere [11], however, it is remarkable that the mechanism is viable for reasonable values of the involved parameters.

IV. CONCLUSIONS

We presented a new scenario where the role of DM is played by heavy RH neutrinos. The scenario is based on a mechanism where the DM RH neutrinos are produced through mixing enhanced by the additional presence of higher dimensional effective operators into the usual type I seesaw Lagrangian. The mechanism relies crucially on the fact that is necessary to convert just a very small fraction of the source RH neutrinos into the DM RH neutrinos. In this way the additional operator has the effect of enhancing the mixing without spoiling any other successful feature of the type I seesaw mechanism and at the same time preserving the DM RH neutrinos stability on cosmological times. A straightforward prediction of the mechanism 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 decaying DM particles with a mass higher than $300\,\mathrm{GeV}$ and a lifetime of approximately $\tau_{DM} \sim 10^{26}\,\mathrm{sec}$ [15]. Therefore, our mechanism seems to have the right features to explain this excess.** These results are quite interesting since not only they are fully compatible with our model but also because the value for the life time corresponds to the saturation of the lower bound Eq. 4 from the AMANDA data when $M_{DM} \sim 10^{5\div 9}$ GeV and a signal should be expected from the ICE CUBE experiment.

It should be also noticed that the special orthogonal form Eq. (9) predicted by the mechanism corresponds [16] to a particular sequential dominated model [17]. Therefore, the proposed scenario for the solution of the DM conundrum restricts remarkably the seesaw parameter space, providing a potential smoking gun for the seesaw mechanism.

We wish to thank F. Bezrukov, S. Davidson, F. Feruglio, A. Ibarra, M. Losada, E. Nardi, S. King, A. Riotto and S. Sibiryakov for useful discussions. PDB is supported by the Helmholtz Association of National Research Centres, under project VH-NG-006.

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[20] In the exact limit $m_1 = 0$ the eigenvalue $h_A = 0$. Moreover, in this limit, plugging $\Omega' = I$ (cf. (Eq. 9)) into the Eq. (7), 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.

[21] The effect of such operator was previosly considered only for nonresonant production of the sterile DM neutrinos (see, e.g., [18]).

[22] This can be also regarded saying that the non-renormalizable operator breaks the symmetry that brings to the special textured form for the Yukawa coupling matrix needed in order to have N_{DM} cosmologically stable.