1. Introduction. The nature of dark matter (DM) is one of the longest-standing puzzles in fundamental physics [1]. Although astrophysical and cosmological observations support a solution in terms of a new particle and disfavor alternative explanations in terms of modified gravity or primordial black holes, all efforts to identify the nature of this new particle, with direct, indirect and collider searches, have failed so far. This produces strong constraints on the existing models of DM, either favoring heavy DM particles (with mass above a TeV), or light particles (with mass below a GeV) or very weakly coupled ones, such as axion-like particles, or some combination thereof. Therefore, a non-thermal production mechanism, quite different from the usual WIMP paradigm, that relies on very small couplings of the DM particle to the thermal bath, is a reasonable possibility to consider.

At the same time extensions of the standard model should also account for neutrino masses and mixing, and explain the matter-antimatter asymmetry of the universe. The type-I seesaw mechanism is the most minimal and attractive way to incorporate neutrino masses and mixing and to explain the matter-antimatter asymmetry of the universe via leptogenesis. It is then quite reasonable to seek unified models of neutrino masses, DM and leptogenesis starting from the type-I seesaw Lagrangian and extended with some new ingredient that also addresses DM.

An example of this kind was proposed in Ref. [2] in which one right-handed (RH) neutrino has vanishing Yukawa couplings and therefore does not contribute to neutrino masses and mixing. However, it mixes with the other RH neutrinos and to the standard model Higgs through the 5-dimensional non renormalizable operator \((\lambda_{11}/\Lambda)\Phi^\dagger \Phi N_{RI} N_{RI} \) [3]. Consequences of these interactions have been studied in Refs. [4,5].

In this Letter we introduce a novel mechanism for the production of RH neutrino dark matter that relies on a first order phase transition of a scalar singlet field. It has been proposed that the variation of couplings in the Weinberg operator during a phase transition can provide a new way to implement leptogenesis [6]. Here we show that a phase transition can induce an efficient conversion of source RH neutrinos into DM RH neutrinos that is compatible with a standard cosmological history. A strong first order phase transition is required and an associated production of gravitational waves (GWs) is expected [7–10]. In fact the spectrum of GWs is related to the properties of the DM particle in a way that provides an interesting signature of the model.

Several works have proposed a more or less direct link of dark matter genesis to a phase transition [7,11–17], some of which also predict a detectable GW spectrum. Additionally, in our scenario there is also an important link with neutrino masses and leptogenesis.

2. The model. In addition to the SM particle content we have three RH neutrinos \(N_{DM}, N_1, N_2\) and a new scalar field \(\eta\) which may be associated with the breaking of a flavor symmetry. The dark RH neutrino \(N_{DM}\) is a neutral gauge singlet which has a Majorana mass \(M_{DM}\) generated by some external mechanism, for example, by symmetry breaking within a dark sector; \(N_{DM}\) plays the role of DM. In a similar way, a source RH neutrino \(N_S\), either \(N_1\) or \(N_2\), acquires a mass by coupling to a different scalar field. The important thing is that the dark sector is weakly coupled to the visible sector through a small mixing of the DM RH neutrino \(N_{DM}\) with \(N_S\) playing the role of a source RH neutrino [2]. This mixing is generated through Higgs portal interactions of \(N_{DM}\) with \(N_S\) and the SM Higgs \(\Phi\), and is described by the operator, \((\lambda_{mix}/\Lambda)\Phi^\dagger \Phi N_{DM}^\dagger N_{DM} N_S\) [2,4]. However, in our case the source RH neutrino \(N_S\) also has a coupling \(\lambda_{S}\) to the new scalar field \(\eta\). Here we do not specify the UV complete model that produces such a mixing but the most attractive option is that the same physics is responsible for both \(\lambda_{mix}\) and \(\lambda_{S}\). Thus, the SM is extended by

\[
\mathcal{L}_\Lambda = \frac{1}{2} M_{DM} N_{DM} \bar{N}_{DM} N_{DM} + \lambda_{S} \eta N_{S}^\dagger N_{S} + \frac{\lambda_{mix}}{\Lambda} \Phi^\dagger \Phi N_{DM}^\dagger N_{DM} N_{S} + h.c.
\]
During the phase transition $\eta$ acquires a vev $v_\eta$ and the source RH neutrino acquires a spacetime-dependent mass.

Including a third RH neutrino, one also recovers the usual type-I seesaw Lagrangian with two RH neutrinos that describes neutrino masses and mixing. At the end of the phase transition the Lagrangian terms extending the SM are (with $I, J = 1, 2$)

$$\mathcal{L}_{Y} = \sum_{\alpha} h_{\alpha I} N_{J} \tilde{\Phi} + \frac{1}{2} M_{I} \mathbb{N}_{I} \delta_{IJ} N_{J} \quad (2)$$

$$+ \frac{\lambda_{\text{mix}}}{\Lambda} \phi^{I} \phi \mathbb{N}_{\text{DM}} \mathbb{N}_{S}$$

$$+ \frac{1}{2} M_{\text{DM}} \mathbb{N}_{\text{DM}} \mathbb{N}_{\text{DM}} + \text{h.c.} ,$$

where either $N_1$ or $N_2$ is identified with the source RH neutrino and the other is neutrino is assumed to have negligible or no mixing with the DM RH neutrino. Then, the third RH neutrino plays no role in DM genesis but, in combination with the source RH neutrino, determines neutrino masses and mixing via a two RH neutrino seesaw mechanism which can potentially generate the matter-antimatter asymmetry via leptogenesis [18].

3. Dark matter genesis. The production of the DM abundance can be calculated by solving the density matrix equation $(I, J = \text{DM, S})$

$$\frac{dN_{IJ}}{dt} = -i [\Delta \mathcal{H}, N]_{IJ} - \left( \begin{array}{c} 0 \\ \Gamma_{\text{dec}} \Gamma_{\text{prod}} \end{array} \right) , \quad (3)$$

where $N_{IJ}$ is the abundance density matrix containing on the diagonal terms the DM and source RH neutrino abundances normalised in a way that in ultra-relativistic thermal equilibrium they are simply unity. The quantities $\Gamma_{\text{dec}}$ and $\Gamma_{\text{prod}}$ are the decoherence and production rates respectively. The effective hamiltonian is given by

$$\Delta \mathcal{H}_{IJ} \simeq \left( \begin{array}{c} -\frac{\Delta \tilde{M}^2}{4 p} \\ \Delta H_{\text{mix}} \end{array} \right) , \quad (4)$$

where $\Delta H_{\text{mix}} = T^2/(12 \Lambda)$, $\Lambda \equiv \Lambda/\lambda_{\text{mix}}$, $\Delta \tilde{M}^2 = \tilde{M}_{S}^2(x, t) - M_{\text{DM}}^2$, with the source RH neutrino effective thermal mass given by the sum of a spacetime dependent term and thermal terms:

$$\tilde{M}_{S}^2(x, t) = M_{S}^2(x, t) + \frac{T^2}{4} h_{S}^2 + \frac{T^2}{8} \lambda_{S}^2 N_{N_{S}} N_{N_{T}} . \quad (5)$$

In this expression $h_{S}^2 \equiv (h^\dagger h)_{SS}$ and we introduce the usual effective neutrino mass, $\tilde{m}_{S} = v^2 h_{S}^2 / M_{S}$. The seesaw mechanism requires $\tilde{m}_{S}/m_{\text{sol}} \geq 1$, where $m_{\text{sol}} \simeq 8.6 \text{ meV}$ is the solar neutrino mass scale. We take the minimum value $\tilde{m}_{S}/m_{\text{sol}} = 1$ that maximizes the DM lifetime, so that $h_{S}^2 = m_{\text{sol}} M_{S}/v^2$.

The spacetime-dependent mass $M_{S}(x, t) = \lambda_{S} v_{\eta}(x, t)$ is generated by the vev of $\eta$ that is described by a well known kink solution for the bubble wall profile [19],

$$v_{\eta}(x, t) = \frac{1}{2} \tilde{v}_{\eta} \left[ 1 + \tanh \left( \frac{x + v_{w}(t - t_{*})}{\Delta_{w}} \right) \right] , \quad (6)$$

where $v_{w}$ and $\Delta_{w}$ are the bubble wall velocity and width, respectively. We comment later on the dependence of our results on the description of the bubble profile.

We assume that before the phase transition the $\eta$ abundance $N_{\eta}$ gets thermalized, so that we can set $N_{\eta} = N_{\eta}^{\text{eq}} = N_{\eta} = 4/3$ in Eq. (3). On the other hand, the source RH neutrino abundance $N_{S}$ is obtained by solving Eq. (3). We adopt a monochromatic approximation with $p \simeq 3T$. The production rate $\Gamma_{\text{prod}}$ for the $N_{S}$ abundance is given by the sum of two contributions: $\Gamma_{D} + \Gamma_{S}$, from inverse decays and scatterings from the Yukawa couplings to the Higgs, and $2 \Gamma_{\eta} \rightarrow N_{S} N_{S}$, from the decays of $\eta$ into source RH neutrinos. Therefore, explicitly we can write

$$\frac{dN_{N_{S}}}{dt} \simeq -(\Gamma_{D} + \Gamma_{S} + 2 \Gamma_{\eta} \rightarrow N_{S} N_{S}) (N_{N_{S}} - N_{N_{S}}^{\text{eq}}) N_{\eta} , \quad (7)$$

where we linearized the term from $\eta$ decays using $N_{N_{S}} \ll N_{N_{S}}^{\text{eq}}$ and also neglected a tiny oscillatory term from the Liouville-von Neumann term from the mixing with $N_{\text{DM}}$.

Introducing the variable $z = T_{*}/T$, where $T_{*}$ is a convenient energy scale that we identify with the temperature of the phase transition, Eq. (7) can be recast as

$$\frac{dN_{N_{S}}}{dz} \simeq -(D + S + D_{\eta}) (N_{N_{S}} - N_{N_{S}}^{\text{eq}}) , \quad (8)$$

where, having introduced the expansion rate,

$$H(z) = \sqrt{\frac{8 \pi^{3} g_{s}}{90} T_{*}^{2}} \frac{1}{z} \zeta^{2} \simeq 1.66 \sqrt{g_{s}} T_{*}^{2} \frac{1}{M_{Pl}} \frac{1}{z} , \quad (9)$$

we defined $(D, S, D_{\eta}) \equiv (\Gamma_{D}, \Gamma_{S}, 2 \Gamma_{\eta} \rightarrow N_{S} N_{S})/(H(z) z)$. For $T \gg M_{S}$ one has $S + D \simeq 2 (\tilde{m}_{S}^{2}/m_{\text{sol}}^{2}) (M_{S}/T_{*})^{2}$. The term from $\eta$ interactions can be written in the form

$$D_{\eta} = \frac{1}{2} \tilde{m}_{S}^{2} (m_{\text{sol}}^{2} T_{*}^{2}) \quad (10)$$

where $m_{\text{sol}} \equiv 16 \pi^{5/2} \sqrt{g_{s}}/(3 \sqrt{5}) (v^{2}/M_{Pl}) \simeq 0.08 \text{ meV}$ is the usual equilibrium neutrino mass and, in a similar fashion to Yukawa couplings, we have also introduced the effective neutrino mass $\tilde{m}_{S} \equiv v^{2} \lambda_{S}^{2}/T_{*}$ that parameterizes the coupling of the source RH neutrino to the scalar $\eta$. The scalar $\eta$ has a vanishing or negligible bare mass but it gets a thermal mass $m_{\eta} = g_{s} T$, where $g_{s}$ is the $\eta$ coupling constant to the thermal bath and typical values are $g_{s} \sim 0.1$. For example, for the standard model Higgs boson, $m_{H}/T \simeq 0.4$. Henceforth, we set $g_{s} = 1/6$. Since in our case the scale $T_{*}$ can be identified with the temperature of the phase transition and it will turn out that $M_{S}/T_{*} \ll 1$, the Yukawa coupling interactions can be neglected in Eq. (7), i.e., $D + S \ll D_{\eta}$. 
Moreover, since the DM is produced when $T \gg M_S$, we approximate $N_{N_0}^{\text{eq}} \simeq 1$ in Eq. (7) finding
\[ N_{N_0}(z) \simeq 1 - e^{-\frac{1}{2} \frac{\tilde{m}_N}{m_*}} . \]  
(11)

We now calculate the dark matter RH neutrino abundance by solving Eq. (4). This can be done by noticing that before the phase transition, $M_S(x, t) = 0$ in Eq. (5). In this way thermal medium effects dominate and suppress the mixing since the off-diagonal term $\Delta H_{\text{mix}}$ is negligible compared to the contribution from $\eta$ interactions in Eq. (5). During the phase transition, the source RH neutrino mass $M_S(x, t)$ increases from zero to its final value $M_S = \lambda_S \bar{v}_\eta r_\eta$. Therefore, during the phase transition, there can be a time when $\Delta \tilde{M}^2(x, t_{\text{res}}) = 0$, or equivalently $\tilde{M}_S^2(x, t_{\text{res}}) = \tilde{M}_{\text{DM}}^2$, corresponding to a mixing resonance condition between $N_S$ and $N_{\text{DM}}$:  
\[ M_{\text{DM}}^2 \simeq M_S^2 \left\{ \frac{1}{4} \left[ 1 + \tanh \left( \frac{x + v_w (t_{\text{res}} - t_*)}{\Delta w} \right) \right]^2 + \frac{1}{6} \left( \frac{T_*}{v_\eta} \right)^2 N_{N_0}(T_*) \right\} , \]  
(12)

where we neglected the effective potential generated by the standard Yukawa coupling term since it is much smaller than the term generated by the $\eta$ interactions. Typically, $T_* \sim \bar{v}_\eta$ and, for definiteness, we fixed $T_* / v_\eta = 1$. The resonance condition (12) shows that, for given values of $T_*$ and $M_S$, there is a finite range $M_{\text{DM}}^2 \leq M_{\text{DM}}^\text{max} \leq M_S^2$ corresponding to a variation of the hyperbolic tangent within $[-1, 1]$. Notice, moreover, that from Eq. (11), because $z_* = 1$, $N_{N_0}(z_*)$ depends only on $\tilde{m}_N^2 = m_*^2 / T_*^2$. Therefore, for a fixed value of $T_*$, the resonance condition (12) identifies an allowed region in the $(M_S, M_{\text{DM}})$ plane. In Fig. 1 we show such regions for three values of $T_*$. Since $T_*$ is so much larger than $M_S$, we have $\tilde{m}_N^2 / m_* \ll 1$, and from (11),  
\[ N_{N_S}(T_*) \simeq \frac{1}{2} g_2^* \frac{\tilde{m}_S^2}{m_*} . \]  
(13)

From Eq. (4), if the resonance is crossed quickly enough in a way that the expansion and the variation of the off-diagonal term $\Delta H_{\text{mix}}(t)$ can be neglected during the crossing, and if the decoherence rate $\Gamma_{\text{dec}} = (1/2) \Gamma_{\text{prod}}$ is negligible compared to $(d \Delta \tilde{M}^2(t)/dt)_{\text{res}} / (6 T_*)$, then the relic DM abundance is very well described by the Landau-Zener formula  
\[ N_{\text{DM}}^{\text{res}} \simeq 12 \pi N_{N_0}(T_*) T_* \frac{\Delta H_{\text{mix}}^2}{d \Delta \tilde{M}^2/dt} \bigg|_{t_{\text{res}}} . \]  
(14)

Notice that a Landau-Zener description was also used in the case of a standard cosmological expansion without a phase transition in Ref. [5]. In that case the bare mass $M_\phi$ is constant, and for $M_\phi \gg M_{\text{DM}}$ the Landau-Zener formula greatly overestimates the DM production. This is because at the resonance, $\Delta \tilde{M}^2(6E) \ll H$, so that RH neutrino oscillations simply do not have time to develop and the Landau-Zener formula breaks down [20].

However, in our case, because of the time dependence of $M_S$ during the phase transition, the value of $\Delta \tilde{M}^2$ at the resonance is sufficiently small that the resonance occurs when oscillations have developed, and we verified that the Landau-Zener formula works quite well, so long as the resonance is crossed quickly enough. One might think of enhancing DM production by making $(d \Delta \tilde{M}^2/dt)_{\text{res}}$ smaller, by taking a smaller $v_w$, but too small a value of $v_w$ invalidates the Landau-Zener formula which relies on a first order Taylor expansion about $t_{\text{res}}$. We take $v_w = 0.9$ and 0.95 to guarantee the validity of the Landau-Zener formula.

We calculate $(d \Delta \tilde{M}^2/dt)_{\text{res}}$, and arrive at the DM abundance produced at the resonance:  
\[ N_{\text{DM}}^{\text{res}} \simeq \frac{N_{N_0}(T_*) (\pi/48) (\Delta w / v_w) T_*^5}{\Lambda^2 M_S^2 (x, t_{\text{res}}) [1 - M_S(x, t_{\text{res}})/M_S]} . \]  
(15)

From this we calculate the present energy density using the relation,  
\[ \Omega_{\text{DM}} h^2 \simeq 1.0875 \times 10^6 N_{\text{DM}}^{\text{res}} M_{\text{DM}} / \text{GeV} . \]  
(16)

This has to be compared with the measured value $\Omega_{\text{DM}} h^2 = 0.11933 \pm 0.00091$ (68\% C.L.) [21] from Planck satellite observations of CMB anisotropies (combining TT,TE,EE +lowE+lensing data sets).

An important point is that even at zero temperature the dark RH neutrino decays because of the mixing with the source RH neutrino. For viable DM masses, which are well below the Higgs mass, the dominant decay mode is $N_{\text{DM}} \rightarrow \nu_{\ell} t_{\ell} \bar{t}_{\ell}$ ($\alpha = e, \mu, \tau$) with a rate,  
\[ \Gamma_{N_{\text{DM}} \rightarrow \nu_{\ell} t_{\ell} \bar{t}_{\ell}} = \frac{\theta^2 N_0}{96 \pi^3 M_S^2} G_F^2 M_{\text{DM}}^2 \]  
(17)
where \( \theta_{A0} = (2 v^2)/(\bar{\Lambda} (M_S - M_{DM})) \) is the mixing angle between the DM and source RH neutrinos at zero temperature \([1]\). The \( \Delta \) decays provide various experimental signatures and tests. For the values of \( M_{DM} \) of interest, the most robust lower bounds on the DM lifetime, \( \tau_{\Delta} \gtrsim 10^{25} \text{s} \), are provided by CMB anisotropy data via changes to the ionization and temperature history \([22]\); see the Planck labelled black dashed curve in Fig. 2, that has been obtained. Data from diffuse X and \( \gamma \)-ray observations also place a lower bound \( \tau_{\Delta} \gtrsim 10^{25} \text{s} \) \([23]\) (not shown in Fig. 2) in the range 0.1 GeV–10 GeV. The points in Fig. 2 correspond to sets of parameter values that reproduce the measured DM abundance. They are obtained by scanning over \( M_{DM}, M_S, \tau_* \) (uniform in logarithm) and \( \tau_{\Delta} \) (instead of \( \Lambda \)). The solutions are divided into five subsets corresponding to the five indicated ranges of \( \tau_* \). There is a tendency for the value of the DM mass to increase with \( \tau_* \), which results from the resonance condition Eq. (12) that, as we have seen from Fig. 1, imposes a relation between \( M_{DM}, M_S \) and \( \tau_* \).

4. Spectrum of gravitational waves. We now estimate the potential for current and future GW interferometers to detect the stochastic GW background produced by the strong first-order cosmological phase transition of \( \eta \). We consider the case of non-runaway bubbles \([24, 25]\), in which the scalar field contribution to GWs is ignored, and conservatively assume that only 5\% of the bulk motion of the bubble walls is converted into vorticity, as supported by numerical results \([26]\). Then, the contribution from sound waves is the dominant source of the GW background in most of the frequency range. The GW spectrum generated by sound waves during the phase transition is determined by four parameters: the phase transition temperature \( T_* \); the bubble wall velocity \( v_w \); the ratio of the vacuum energy density released in the transition to the vacuum energy density of the radiation bath \( \alpha \); the fraction \( \beta/H_* \), where \( H_* \) is the expansion rate at \( T_* \) and \( \beta \) is (approximately) the inverse of the time duration of the phase transition. The first two parameters are irrelevant for DM genesis. We only consider values of \( \alpha \) that permit \( v_w \) much larger than \( c_s(1+\sqrt{2} \alpha) \), where \( c_s = 1/\sqrt{3} \) is the sound speed in the plasma \([25]\). In this regime the analytical results and fits provided in Ref. \([24]\) can be employed.

We display the GW spectrum for our five benchmark points in Fig. 3. Sensitivities of the LIGO O2 & O5 observing runs \([27]\), LISA \([24, 28]\), ET \([29]\), BBO \([30]\), and DECIGO \([31]\) are shown for comparison. TianQin \([32]\), and Taiji \([33]\) have sensitivities similar to LISA. As shown in Fig. 2, successful DM genesis requires \( T_* \gtrsim 10^3 \) GeV for the bulk of the points. LISA can test this regime. LIGO, ET, and BBO might be able to test a high temperature phase transition, \( T_* \sim 10^7 \text{ GeV} – 10^8 \text{ GeV} \), if \( \alpha \) takes a relatively large value. This is quite interesting because then \( M_S \) can be sufficiently large \( (M_S \gtrsim 300 \text{ GeV}) \) to make resonant leptogenesis viable.

4. Final remarks. We introduced a novel model in which DM and neutrino masses are generated during a phase transition of a scalar field. The phase transition needs to be strongly first order for efficient DM genesis, which has the welcome feature of the associated production of GWs with an intensity measurable at the planned GW interferometers, LISA, BBO, DECIGO, ET, TianQin and Taiji. We find solutions for DM masses within a range 1 MeV–20 GeV with a clear tendency of the temperature of the phase transition to increase with the dark matter mass, though the dependence also involves other parameters such as the mass of the source RH neutrino. Interestingly, for high values of the phase transition temperature, \( T_* \sim 10^7 \text{ GeV} \), solutions with source RH neutrinos heavier than 300 GeV are possible which also allows successful resonant leptogenesis \([34]\) with two RH neutrinos \([5]\). For masses \( \sim \text{GeV} \), leptogenesis from RH neutrino oscillations \([35]\) is a possibility. We have also seen that decays of the DM neutrino decays offer an ad-

| \( T_* \) | \( \tau_{\Delta} \) | \( M_S \) | \( M_{DM} \) | \( v_w \) | \( \alpha \) | \( \beta/H_* \) |
|---|---|---|---|---|---|---|
| B1 | 0.0033 | 0.37 | 5.35 | 1.80 | 0.90 | 0.10 | 10 |
| B2 | 0.0635 | 1.78 | 133 | 19.73 | 0.90 | 0.10 | 10 |
| B3 | 0.486 | 1.08 | 805 | 37.36 | 0.90 | 0.10 | 10 |
| B4 | 6.64 | 3.80 | 6.4 \times 10^3 | 176.1 | 0.95 | 0.15 | 5 |
| B5 | 37.7 | 0.74 | 411 \times 10^3 | 15.4 \times 10^3 | 0.95 | 0.15 | 5 |

**TABLE I:** Benchmark points in Fig. 2 obtained for \( v_w/\Delta w = T_*/10 \).
ditional way to test the scenario through observations of CMB anisotropies and/or the diffuse X- and γ-ray backgrounds. We have not thoroughly studied the dependence of our results on all parameters and have fixed some of them to reasonable values, e.g., $g_0 = 1/6$, $v_w = T$, and $\tilde{m}_S/m_{\text{rad}} = 1$. However, we verified that our results are insensitive to small variations of the first two parameters, and that increasing $\tilde{m}_S$ reduces the allowed range of DM masses. Consequently, since $\tilde{m}_S \gtrsim m_{\text{atm}}$ for the inverted neutrino mass hierarchy, a significant reduction of the allowed parameter space ensues, which implies that our scenario clearly prefers a normal mass hierarchy. We also assumed a specific bubble wall profile with a fixed value of the bubble width and velocity. A thorough exploration of the allowed values of all parameters will be presented in a forthcoming paper together with the possibility of some variants of the model and a discussion of different aspects, including whether a contribution from magneto hydrodynamic turbulence might give a sizable contribution that enhances the signal from primordial GWs [30]. However, our results clearly prove the viability of the proposed mechanism and how DM genesis, although escaping traditional direct and collider searches, will be tested in the future with cosmology, cosmic rays and GW interferometers.

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