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Sterile neutrino dark matter and low scale leptogenesis from a charged scalar

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Abstract We show that novel paths to dark matter generation and baryogenesis are open when the standard model is extended with three sterile neutrinos N_i and a charged scalar δ^+. Specifically, we propose a new production mechanism for the dark matter particle—a multi-keV sterile neutrino, N_1—that does not depend on the active-sterile mixing angle and does not rely on a large primordial lepton asymmetry. Instead, N_1 is produced, via freeze-in, by the decays of δ^+ while it is in equilibrium in the early Universe. In addition, we demonstrate that, thanks to the couplings between the heavier sterile neutrinos $N_{2,3}$ and δ^+, baryogenesis via leptogenesis can be realized close to the electroweak scale. The lepton asymmetry is generated either by $N_{2,3}$-decays for masses $M_{2,3} \gtrsim$ TeV, or by $N_{2,3}$-oscillations for $M_{2,3} \sim$ GeV. Experimental signatures of this scenario include an X-ray line from dark matter decays, and the direct production of δ^+ at the LHC. This model thus describes a minimal, testable scenario for neutrino masses, the baryon asymmetry, and dark matter.

1 Introduction

Sterile neutrinos, that is, fermions singlet under the $SU(3) \times SU(2) \times U(1)$ gauge symmetry, are a very well-motivated extension of the standard model (SM). On the theoretical side, they are a prediction of left-right symmetric theories, they allow to gauge $B - L$ by removing its anomaly and they are necessary in $SO(10)$ grand unification. On the phenomenological side, they provide a non-vanishing mass to the active SM neutrinos, they allow to realize baryogenesis via leptogenesis, and they are also a viable candidate for dark matter, as long as their mass lies in the keV range. In that case, they are naturally long-lived so that, unlike dark matter candidates at the electroweak scale, no additional symmetries are required to stabilize them. Thus, sterile neutrinos

can provide a simple solution to the three open problems of the SM: neutrino masses, the baryon asymmetry, and the dark matter.

The minimal model addressing these three issues requires one sterile neutrino N_1 at the keV scale as dark matter candidate [1], and two additional sterile neutrinos $N_{2,3}$ for leptogenesis, which is induced either by N-decays, for sterile neutrino masses above the TeV scale [2], or by N-oscillations, for sterile neutrino masses at the GeV scale [3,4]. Light active neutrino masses are easily generated, via the seesaw mechanism, provided the sterile neutrino masses are significantly larger than about 1 eV. Thus, one may argue that no new physics is needed above the electroweak scale to explain neutrino masses, baryogenesis, and dark matter, defining the so-called "ν minimal standard model" (νMSM) [5]. Such a model is undoubtedly economical and very predictive, by reason of the small number of degrees of freedom it contains, but it is also strongly constrained. Present experiments, in fact, already exclude the minimal scenario for sterile neutrino dark matter within the νMSM [6]. In the region of mass and mixing angle compatible with current observations, in particular the X-ray bounds [7–9] and the Ly$_\alpha$ forest data [10–12], one cannot produce a large enough population of sterile neutrinos from standard active-sterile oscillations [1]. The only way out within the νMSM is to generate the dark matter neutrinos via resonant active-sterile oscillations triggered by large initial lepton asymmetries [13,14], but that mechanism requires a high degree of fine-tuning [6,15]. In addition, leptogenesis via $N_{2,3}$-oscillations within the νMSM also requires a significant tuning of parameters, in particular a strong mass degeneracy between N_2 and N_3 [6].

It is important, therefore, to consider alternative ways of realizing leptogenesis and producing sterile neutrinos within extensions of the νMSM. Indeed, most extensions of the SM that are theoretically well-motivated contain new degrees of freedom beside the sterile neutrinos, which may modify significantly the phenomenology of the νMSM. For leptogen-

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esis, the required tuning of parameters can be released by introducing a second Higgs doublet, as recently studied in [16]. Regarding dark matter, several other possibilities have been considered for the production of the sterile neutrinos. They include the production through inflaton decay [17–19], thermal overproduction followed by entropy dilution [20], and the production via the decays of either a neutral scalar in thermal equilibrium [21], or a frozen-in neutral scalar [22].

In this paper we advocate an extension of the SM by three sterile neutrinos, N_i, plus a charged scalar, δ^+, both of which are naturally present in left-right symmetric or unified extensions of the SM [23]. The charged scalar interacts with SM leptons, both doublets and singlets, and with sterile neutrinos, and its mass can be as low as allowed by current collider bounds—about 200 GeV. We will show that the interplay between N_i and δ^+ modifies the dark matter phenomenology, enables new realizations of baryogenesis, and gives rise to novel experimental signatures at colliders. The decays of δ^+, in fact, constitute a new mechanism for the production of sterile neutrino dark matter in the early Universe. An advantage of this mechanism is that the resulting relic density does not depend on the active-sterile mixing angle, allowing to satisfy the strong bounds derived from X-ray observations [7–9]. Besides, the charged scalar induces independent sources of leptogenesis that are effective at scales close to the electroweak scale. If the charged scalar mass were accessible at colliders, the present scenario for dark matter and baryogenesis would be testable in a direct way.

The rest of the paper is organized as follows. In Sect. 2 we introduce the model and fix our notation. Section 3 deals with dark matter production via the decays of δ^+. We obtain, in particular, the dark matter relic density as a function of the parameters of the model. Baryogenesis via leptogenesis is discussed in Sect. 4. We examine two different ways to generate a lepton asymmetry, and outline the region of the parameter space that successfully realizes baryogenesis in each case. The most interesting experimental signatures of this model are discussed in Sect. 5. Section 6 briefly elaborates on the flavour structure of the masses and Yukawa couplings of the model. We present our conclusions in Sect. 7.

2 The model

When gauge singlet chiral fermions N_{Ri} are added to the SM, they have in general a Majorana mass term and a Yukawa coupling to the SM lepton doublets,

$$
\mathscr{L}_N = \overline{N_{Ri}} i \gamma^\mu \partial_\mu N_{Ri}
$$
$$
+ \left[-\frac{1}{2} \overline{(N_{Ri})^c} (M_N)_{ij} N_{Rj} - \overline{l_{L\alpha}} (y_\nu)_{\alpha i} N_{Ri} \tilde{H} + h.c. \right].
\tag{1}
$$

When a scalar field δ^+, singlet under $SU(2)$ and with electromagnetic charge one, is added to the SM, it has in general, besides its gauge and self-interactions, a quartic coupling to the Higgs doublet and an antisymmetric Yukawa coupling to two lepton doublets,

$$
\mathscr{L}_\delta = D_\mu \delta^+ D^\mu \delta^- - M_\delta^2 \delta^+ \delta^- - \frac{1}{2} \lambda_\delta (\delta^+ \delta^-)^2
$$
$$
- \lambda_{\delta H} \delta^+ \delta^- H^\dagger H + \left[-\overline{l_{L\alpha}} (y_L)_{\alpha\beta} (i\sigma_2) (l_{L\beta})^c \delta^- + h.c. \right].
\tag{2}
$$

When both the sterile neutrinos N_{Ri} and the charged scalar δ^+ are present, there is an additional Yukawa coupling involving the SM lepton singlets,

$$
\mathscr{L}_{\delta N} = -\overline{(e_{R\alpha})^c} (y_R)_{\alpha i} N_{Ri} \delta^+ + h.c.
\tag{3}
$$

These three terms, Eqs. (1)–(3), plus the well-known SM ones, constitute the Lagrangian considered in this paper. It is the most general Lagrangian consistent with the $SU(3) \times SU(2) \times U(1)$ gauge symmetry and with the addition of the singlet fermions and the charged scalar. Without loss of generality, we adopt the basis where M_N and the charged lepton mass matrix are real and diagonal. For our following discussion, the parameters of greater relevance are the masses of the singlet fermions (M_i, $i = 1, 2, 3$) and of the charged scalar (M_δ), and the 3×3 Yukawa matrices y_ν and y_R.

In this model active neutrino masses are not affected by the existence of δ^+ and are obtained via the usual seesaw mechanism: $m_\nu \simeq y_\nu M_N^{-1} y_\nu^T v^2$, with $v \simeq 174$ GeV. In Sect. 3 we will identify the sterile neutrino N_1 with a multi-keV dark matter candidate, with very small mixing angles with active neutrinos, corresponding to tiny neutrino Yukawa couplings, $(y_\nu)_{\alpha 1} \lesssim 3 \times 10^{-13}$. These tiny couplings give rise to a lightest neutrino mass $m_\nu^{\text{lightest}} \lesssim 10^{-6}$ eV. The much larger solar and atmospheric neutrino mass scales are generated instead by the seesaw contribution of the two heavier sterile neutrinos, $N_{2,3}$. The active neutrino mass spectrum is thus hierarchical (either normal or inverted), with one neutrino remaining essentially massless.

The generation of the dark matter energy density and of the baryon asymmetry have been extensively studied in the context of the SM extended with sterile neutrinos only [6]. The SM extension with sterile neutrinos and δ^+ was considered in [23]. It was shown that, taking the sterile neutrinos N_i as light as a few TeVs and heavier than δ^+, one can generate the observed baryon asymmetry via leptogenesis, with no need to enhance resonantly the CP asymmetry.

We explore, instead, the possibility that the lightest singlet fermion, N_1, has a multi-keV mass, $M_1 \ll M_\delta$, and accounts for the observed dark matter density. In this framework, the existence of δ^+ offers an alternative way of producing dark matter that has not been studied before in the

literature. Moreover, we will show that this different mass spectrum is still compatible with the leptogenesis mechanism proposed in [23] and, in addition, we will identify a different mechanism for leptogenesis in the presence of δ^+ based on $N_{2,3}$-oscillations.

3 Dark matter production from δ^+ decays

In this model, the only particle that can play the role of dark matter is the lightest singlet fermion, N_1. It is important, therefore, to determine how it is produced in the early Universe and whether it can account for the observed dark matter density while respecting all other experimental constraints. In this section, after briefly reviewing the standard scenario for sterile neutrino dark matter, we propose a new mechanism for dark matter production via the decays of δ^+.

In the absence of δ^+, our model is described by the seesaw Lagrangian, Eq. (1), and the region of the parameter space compatible with dark matter is the one of the νMSM. It features a singlet fermion (N_1) with a mass at the keV scale, the dark matter particle, and two heavier singlets, $N_{2,3}$. In the νMSM, dark matter is produced at temperatures of order 100 MeV via active-sterile neutrino mixing. The required effective mixing angle, θ_1, is necessarily small and is related to the Yukawa couplings by $\theta_1^2 = \sum_{\alpha=e,\mu,\tau} |(y_\nu)_{\alpha 1}|^2 v^2/M_1^2$. Detailed numerical studies have shown that the observed dark matter density can be reproduced in two different cases [6]. In the non-resonant case (also known as the Dodelson–Widrow mechanism [1]), dark matter neutrinos are thermally and non-resonantly produced with a smooth distribution of momenta. This mechanism yields the minimal amount of dark matter that can be obtained for a given mass and mixing angle. For dark matter masses between 1 keV and 10 keV, the mixing angle required to explain the dark matter density is $\sin^2(2\theta_1) \sim 10^{-8}$-$10^{-9}$, as shown e.g. in figure 2 of [6]. These parameters are not consistent with current data [24], which imply $M_1 \lesssim 3-4$ keV from the X-ray line bound [7–9] and $M_1 \gtrsim 8$ keV from the Ly$_\alpha$ forest observations [10–12]. The other case in which the dark matter density can be explained, and the only one that is currently viable in the νMSM, is resonant production (also known as the Shi-Fuller mechanism [13]). In it, the dark matter production rate is resonantly amplified by the presence of a lepton chemical potential in the plasma, which enhances the production for particular momenta as they pass through the resonance, giving rise to a non-thermal momentum distribution that is *colder* than that obtained in the non-resonant case. To explain the observed dark matter density, a lepton asymmetry $|\mu_\alpha| \gtrsim 8 \times 10^{-6}$ at $T \sim 100$ MeV is required [14], where $\mu_\alpha = n_\alpha/s$, s is the entropy density of the Universe, and n_α is the total number density of active leptons of flavour α. In principle, this large lepton asymmetry can be generated

within the νMSM via CP-violating oscillations of N_2 and N_3, but only if their mass difference, $M_3 - M_2$, is fine-tuned to the order of 10^{-11} [15,25].

The existence of δ^+ allows a new production mechanism for sterile neutrinos in the early Universe. Indeed, thanks to the interaction term in Eq. (3), the decays of the charged scalars can produce sterile neutrinos, $\delta^+ \rightarrow N_1 \ell^+$. This decay is a thermal process that takes place while δ^+ is in equilibrium and it is an example of the so-called freeze-in scenario [26] for dark matter production. The defining feature of freeze-in is that the dark matter interactions are so weak that they never reach thermal equilibrium in the early Universe. As a result, the dark matter abundance is negligible at high temperatures and slowly increases as the Universe cools down. The production ceases when the freeze-in temperature is reached, and the dark matter abundance remains constant from then on.

The N_1 yield, $Y_{N_1}(T) = n_{N_1}(T)/s(T)$, from δ^+ decays is obtained by solving the following Boltzmann equation [26]

$$sT\frac{dY_{N_1}}{dT} = -\frac{\gamma_{N_1}(T)}{H(T)} \tag{4}$$

where s is the entropy density, H is the expansion rate and $\gamma_{N_1}(T)$ is the thermally-averaged production rate. We have that

$$\gamma_{N_1}(T) = \frac{M_\delta^2 T}{2\pi^2} K_1(M_\delta/T)$$
$$\times \sum_\alpha \left[\Gamma(\delta^- \rightarrow N_1\ell_\alpha^-) + \Gamma(\delta^+ \rightarrow N_1\ell_\alpha^+)\right] \tag{5}$$

where $K_1(x)$ is the Bessel function of the second kind and the sum runs over the different lepton flavours. Since the lepton and N_1 masses are negligible compared to M_δ, the δ^+ decay rates are calculated as

$$\Gamma(\delta^- \rightarrow N_1\ell_\alpha^-) = \Gamma(\delta^+ \rightarrow N_1\ell_\alpha^+) = \frac{M_\delta|(y_R)_{\alpha 1}|^2}{16\pi} \tag{6}$$

and their sum over flavours can be conveniently written as

$$\sum_\alpha \left[\Gamma(\delta^- \rightarrow N_1\ell_\alpha^-) + \Gamma(\delta^+ \rightarrow N_1\ell_\alpha^+)\right] = \frac{M_\delta}{8\pi}\left(y_R^\dagger y_R\right)_{11}$$
$$= \frac{M_\delta}{8\pi}y_{R1}^2, \tag{7}$$

where we have defined y_{R1} as the combination of couplings that determines the production of N_1 via δ^+ decays. Using $s(T) = 2\pi^2 g_s T^3/45$, $H(T) = 1.66\sqrt{g_\rho}T^2/M_{Pl}$, and $x \equiv M_\delta/T$ we can solve Eq. (4) as

$$Y_{N_1}(T_0) = \frac{45}{(1.66)32\pi^5 g_s \sqrt{g_\rho}}\frac{M_{Pl}y_{R1}^2}{M_\delta}\int_{x_{min}}^{x_{max}} x^3 K_1(x)\,dx, \tag{8}$$

where T_0 is the temperature today and we used $Y(T \gg M_\delta) = 0$ as our initial condition. Integrating this equation from $x_{max} = \infty$ to $x_{min} = 0$ and setting $g_s \approx g_\rho \approx 100$ yields

$$Y_{N_1}(T_0) \approx 1.6 \times 10^{-5} \left(\frac{y_{R1}}{10^{-8}} \right)^2 \left(\frac{1 \text{ TeV}}{M_\delta} \right). \tag{9}$$

The N_1 relic density is then

$$\Omega_{N_1} h^2 \approx 0.11 \left(\frac{M_1}{\text{keV}} \right) \left(\frac{y_{R1}}{5 \times 10^{-8}} \right)^2 \left(\frac{1 \text{ TeV}}{M_\delta} \right). \tag{10}$$

Thus, a keV sterile neutrino can explain the observed dark matter density via δ^+ decays if $M_\delta \sim 1$ TeV and $y_{R1} \sim$ few $\times 10^{-8}$. This analytical result is in very good agreement with the numerical calculations we did for similar models of freeze-in dark matter [27,28].

A crucial feature of Eq. (10) is that, unlike the production in the νMSM, it does not depend on the active-sterile mixing angle θ_1 that determines the decay rate of the sterile neutrino. Hence, one can take θ_1 small enough to be in agreement with the X-ray bounds without affecting the predicted dark matter density. In other words, in this model one can decouple the dark matter production (determined by the y_R couplings) from the dark matter decay (determined by y_ν via mixing angles), opening new viable regions and alleviating the experimental constraints on sterile neutrino dark matter. In particular, M_1 can be larger than the keV scale, as long as the dark matter sterile neutrino is sufficiently long-lived. Since $\Gamma_{N_1} \propto M_1^5 \sin^2 \theta_1$, one cannot exclude, for example, the possibility of vanishing active-sterile mixing, $\theta_1 = 0$, which would allow masses as high as 1 MeV. Above that value, the decay $N_1 \to \nu e^+ e^-$ becomes kinematically available and can proceed through the couplings y_R and y_L, even for $\theta_1 = 0$. In that case, the strong constraints coming from indirect searches of electrons and positrons come into play, but they compete with the smallness of y_{R1}, determined by Eq. (10), and of the Yukawa matrix y_L, that can be very small too.

Regarding structure formation, the sterile neutrinos produced via δ^+ decays are *colder* than those obtained in the νMSM [29,30] because the decays take place at a temperature $T \sim M_\delta$ much higher than the QCD scale (~ 150 MeV). In fact, their free-streaming length and phase space density are identical to those of the so-called chilled sterile neutrinos studied in [31]. Consequently, the Lyman-α bounds on the sterile neutrino mass in our scenario are significantly weaker, allowing for a dark matter mass M_1 smaller than in the νMSM [32].

We assume in the following that the decays of the charged scalars are the dominant source of sterile neutrinos, so that Eq. (10) is satisfied. This will be the case in the absence

of a large lepton asymmetry and for small mixing angles, $\sin^2 2\theta_1 < 10^{-9}$. Next, we show that this assumption is consistent with baryogenesis, and that it leads to new experimental signatures.

4 Baryogenesis via leptogenesis with δ^+

In this section we will discuss leptogenesis in the presence of the charged scalar field δ^+. We will briefly review the two basic mechanism for leptogenesis in the minimal scenario with sterile neutrinos only: N-decays and N-oscillations. We will show that both mechanisms are still operative when one replaces the Yukawa coupling y_ν in Eq. (1) with y_R in Eq. (3), that is, when the role of the SM Higgs doublet H is played by δ^+, and the role of the SM lepton doublet l_L is played by the SM lepton singlet e_R.

The baryon asymmetry of the Universe can be generated from a lepton asymmetry, as long as the latter is present before the electroweak phase transition, which occurs at a temperature $T_{EW} \simeq 150$ GeV. Above this temperature the $(B + L)$-violating electroweak sphalerons are in thermal equilibrium, thus converting efficiently leptons into baryons. When the SM is extended by sterile neutrinos, one can generate a lepton asymmetry, provided that the set of couplings $(M_N)_{ij}$, $(y_\nu)_{\alpha i}$ and $(y_R)_{\alpha i}$ violates the CP symmetry, and (some of) these couplings are out-of-equilibrium at some epoch before T_{EW}.

4.1 Leptogenesis from N-decays

The traditional leptogenesis mechanism [2] assumes a Majorana mass matrix M_N for two (or more) sterile neutrinos, with eigenvalues M_i larger than the electroweak scale, and it relies on the out-of-equilibrium decays $N_i \to H l_{L\alpha}$, at temperatures just below M_i. As N_1 plays the role of keV-scale dark matter candidate, the relevant decays are those of N_2, with $T_{EW} < M_2 < M_3$. Since M_2 violates lepton number, the CP-asymmetry that is generated in the decays is also a lepton asymmetry, or equivalently a $(B - L)$-asymmetry. After all the N_2 particles decayed, the produced $(B - L)$-asymmetry remains constant. In this scenario, the couplings that control leptogenesis are the same Yukawa couplings, y_ν, that determine the active neutrino masses. As a consequence, a sufficient amount of baryon asymmetry can be generated only for $M_2 \gtrsim 10^8$ GeV approximately.[1] Thus, no direct test of this scenario can be envisaged.

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[1] This lower bound on the leptogenesis scale was pointed out in [33] and its exact value depends on various parameters, in particular the initial N_1 abundance [34]. It can be evaded if one allows for resonance effects, that require a strong degeneracy among the sterile neutrino mass eigenstates [35].

In the presence of a charged scalar δ^+, an analog source of leptogenesis is provided by the decays $N_i \to \delta^+ e_{R\alpha}$, mediated by the Yukawa coupling y_R, as shown in [23]. Here we adapt the results of that paper to our new context, where the dark matter particle is the keV-scale sterile neutrino N_1. We can safely neglect the presence of N_1 for the computation of the lepton asymmetry because its interaction rates (in particular those violating lepton number) are strongly out-of-equilibrium above T_{EW}, due to the extreme smallness of M_1, $(y_R)_{\alpha 1}$ and $(y_\nu)_{\alpha 1}$. Thus, at least two extra sterile neutrinos $N_{2,3}$ heavier than δ^+ are needed to generate the CP-asymmetry. Taking for simplicity $M_3 \gg M_2 \gg M_\delta$, the CP-asymmetry in N_2-decays is given by

$$\epsilon_{N_2} = \frac{1}{8\pi} \frac{\text{Im}[\sum_\alpha (y_R)_{\alpha 2}(y_R)^*_{\alpha 3}]^2}{\sum_\alpha (y_R)_{\alpha 2}(y_R)^*_{\alpha 2}} \frac{M_2}{M_3}, \tag{11}$$

where we assumed that the charged scalar lepton number is $L(\delta^+) = -2$, since the dominant decay mode is $\delta^+ \to e^+ \bar{\nu}$, through the Yukawa coupling y_L in Eq. (2).[2] To reproduce the observed baryon asymmetry two basic conditions are required [23]: a minimal value for ϵ_{N_2} assuming no washout,

$$\frac{|\text{Im}[\sum_\alpha (y_R)_{\alpha 2}(y_R)^*_{\alpha 3}]^2|}{\sum_\alpha (y_R)_{\alpha 2}(y_R)^*_{\alpha 2}} \gtrsim 2 \cdot 10^{-6} \frac{M_3}{M_2}, \tag{12}$$

and N_2-decays out-of-equilibrium at $T = M_2$ to avoid large washout from inverse decays,

$$\sum_\alpha (y_R)_{\alpha 2}(y_R)^*_{\alpha 2} \lesssim 10^{-13} \frac{M_2}{1 \text{ TeV}}. \tag{13}$$

Thus, one can realize leptogenesis for M_2 as small as a few TeVs, as long as $N_{2,3}$ have hierarchical Yukawa couplings: $|(y_R)_{\alpha 2}| \lesssim 3 \times 10^{-7}$ and $|(y_R)_{\alpha 3}| \gtrsim 10^{-3}\sqrt{M_3/M_2}$. It is worth reminding that, when M_2 approaches the electroweak scale, a sufficiently large CP-asymmetry implies in general strong $(B-L)$-washout rates, in particular those mediated by N_3 off-shell, unless special conditions are realized. A neat, model-independent discussion of the lower bound on the leptogenesis scale can be found in [36,37]. In the present model with $M_2 \sim$ a few TeVs, the washout can be Boltzmann suppressed by raising M_δ sufficiently close to M_2, and taking a sufficiently large coupling y_L in Eq. (2) to avoid washout from the asymmetry stored in δ^+.[3]

Note that the asymmetries generated by N_2-decays through the Yukawa coupling matrix y_ν become negligible in the

region $M_2 \ll 10^8$ GeV due to the seesaw relation. The y_R-entries, on the other hand, are not constrained by the light neutrino masses, and leptogenesis can work close to the electroweak scale as described above. This scenario is, therefore, easier to test directly, by the observation of δ^+ (and possibly N_2) at colliders, as discussed in Sect. 5.

4.2 Leptogenesis from N-oscillations

A complementary mechanism for leptogenesis [3] relies on N-oscillations rather than on N-decays. It also requires two (or more) sterile neutrinos, coupled to the SM through the Yukawa matrix y_ν. In the early Universe, the thermal population of lepton doublets $l_{L\alpha}$ produces, through small, out-of-equilibrium couplings $(y_\nu)_{\alpha i}$, a coherent superposition N_α of the sterile neutrino mass eigenstates N_i. These sterile neutrinos, at temperatures much larger than their masses, coherently oscillate among the different flavours α. Such oscillations conserve lepton number (the N_α conserve their helicity), but violate lepton flavour numbers. If the CP-symmetry is also violated, one generates non-zero flavour asymmetries between the opposite helicities of the N_α. The asymmetry in the flavour α is transferred efficiently to $l_{L\alpha}$, as long as the coupling $(y_\nu)_{\alpha i}$ goes into equilibrium for some i. Since the total lepton asymmetry, that is, the sum over α of the flavour asymmetries vanishes, one needs that some but not all flavours go into equilibrium before T_{EW}, so that a net lepton asymmetry remains stored in the sterile neutrino sector, and an opposite one is available in the SM sector to be transferred to the baryons by electroweak sphalerons. The Yukawa (out-of-)equilibrium condition at T_{EW} reads

$$|(y_\nu)_{\alpha i}| \gtrsim 10^{-7} \quad \text{for some } \alpha \text{ and } i, \quad |(y_\nu)_{\beta i}|$$
$$\lesssim 10^{-7} \quad \text{for some } \beta \neq \alpha, \, \forall i. \tag{14}$$

A number of additional constraints, on the size of y_ν-entries and on the values of M_i, must be satisfied for this leptogenesis mechanism to work. To facilitate the comparison between the scenarios with and without the charged scalar, we will describe them in some detail. Due to the dark matter constraint, N_1 plays no role in leptogenesis and the asymmetry must be generated by N_2 and N_3 only.[4]

Successful leptogenesis implies a few upper bounds on the y_ν-entries, beside the flavour-dependent one in Eq. (14). First of all, note that the Majorana masses M_i violate lepton number. Therefore, they play the role of washout in this scenario, as they transform opposite helicities of the N_α into one another. The condition to keep the lepton number violation rate out-of-equilibrium down to T_{EW} is approximately

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[2] The asymmetry changes by an order one factor if y_L is negligibly small and δ^+ undergoes slower three-body decays [23].

[3] To reduce the washout, one can also delay N_2-decays by taking $(y_R)_{\alpha 2}$ very small, but in this case one may need to produce the initial thermal density of N_2 by some other interaction.

[4] Note that the seesaw lagrangian can violate CP with only two sterile neutrinos, as it contains three physical phases, one combination of them being relevant in oscillations.

$$|(y_\nu)_{\alpha i}| \lesssim 10^{-5} \frac{\text{GeV}}{M_i}, \quad \forall \alpha \text{ and } \forall i, \tag{15}$$

which combined with Eq. (14) implies sterile neutrino masses below the electroweak scale, $M_i \lesssim 100\,\text{GeV}$.[5] In addition, for the case of two sterile neutrinos $N_{2,3}$, analytic and numerical studies [4,16,39] show that a sufficient baryon asymmetry requires a strong degeneracy between M_2 and M_3, with $\Delta M/M \lesssim 10^{-5}$. This comes from the interplay of a number of subtle effects: the oscillation time increases as ΔM decreases, and this allows for larger asymmetries because the Yukawa interaction rates are closer to equilibrium at later times; note also that larger Yukawa couplings (for some but not all the flavours) enhance the flavour asymmetries, but also tend to spoil the coherence of the sterile neutrinos, as active-sterile transitions may become faster than the oscillation time. We extrapolate the resulting constraint from figure 7 of [6], that can be written as

$$|(y_\nu)_{\alpha i}| \lesssim 2 \times 10^{-6} \left(\frac{M_i}{\text{GeV}}\right)^{1/2} \quad \forall \alpha \text{ and } \forall i. \tag{16}$$

The masses $M_{2,3}$ are below collider energies, but it is difficult to produce them directly since they have small couplings. Still, the mixing with active neutrinos can be sufficiently large to have an observable effect in various neutrino experiments [40]. The resulting upper bound on active-sterile mixing is also shown in figure 7 of [6]. An experimental proposal to improve significantly the present bound can be found in [41]. In less minimal models, e.g. with three GeV scale sterile neutrinos, there are good detection perspectives in meson decay experiments [42].

On the other hand, several observables other than the baryon asymmetry put lower bounds on the y_ν-entries, that add to the one in Eq. (14). The lower bound on active neutrino masses, $|(m_\nu)_{\alpha\beta}| \gtrsim 0.025\,\text{eV}$ for some α and β, implies a lower bound on the neutrino Yukawa couplings through the seesaw formula,

$$|(y_\nu)_{\alpha i}| \gtrsim 2 \times 10^{-8} \left(\frac{M_i}{\text{GeV}}\right)^{1/2} \quad \text{for some } \alpha \text{ and } i. \tag{17}$$

Indeed, since N_1 gives a negligible contribution to m_ν, one active neutrino is approximately massless, therefore $|(m_\nu)_{\alpha\beta}| \lesssim 0.05\,\text{eV}$ for all α and β. Then, the right-hand

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[5] We remark that most numerical studies in the literature assume that the lepton number violating rates mediated by M_i are negligible; in view of the size of the Yukawa couplings relevant for leptogenesis, this assumption seems to be justified only for $M_i \lesssim \text{GeV}$. For larger M_i, one should include the lepton number violating rates in the Boltzmann equations. In this regime, it remains conceivable that the CP-asymmetries are large enough to compensate the relatively strong washout. This possibility was recently entertained to make this scenario work even for M_i larger than T_{EW} [38].

side of Eq. (17) provides the natural value of the Yukawa couplings for $i = 2, 3$: the largest $(y_\nu)_{\alpha i}$ are, the strongest is the cancellation needed among the contributions of N_2 and N_3 to m_ν. These GeV-scale sterile neutrinos should decay (e.g. in 3ν's or $\nu e^+ e^-$) before $T \simeq 1\,\text{MeV}$, not to spoil nucleosynthesis (see [43] for a detailed analysis). We roughly estimate this constraint as

$$|(y_\nu)_{\alpha i}| \gtrsim 3 \cdot 10^{-8} \left(\frac{\text{GeV}}{M_i}\right)^{3/2}, \quad \text{for some } \alpha, \quad \forall i. \tag{18}$$

Comparing with Eq. (16), this implies $M_i \gtrsim 0.1\,\text{GeV}$. All the constraints above confirm that N_1 plays no role in leptogenesis.

Let us show that the same mechanism of leptogenesis through N-oscillations is operative in the presence of the charge scalar δ^+, by replacing the role of y_ν with y_R. Analogously to the previous case, the couplings $(y_R)_{\alpha i}$ must be small to remain out-of-equilibrium while they slowly produce coherent sterile neutrino states. The lepton flavour asymmetries generated by N-oscillations are (partially) transferred to the SM lepton singlets $e_{R\alpha}$. The latter are in equilibrium with $l_{L\alpha}$ through the charged lepton Yukawa couplings (at least for $\alpha = \mu, \tau$), thus electroweak sphalerons transfer the asymmetries to baryons as usual.

One needs that the y_R-entries satisfy the same inequalities as the y_ν-entries in Eqs. (14) and (15), so that some flavour asymmetries are transferred to baryons, while the others remain stored in the sterile sector, and the washout from lepton number violating scattering is small. The only differences amount to (i) order one factors to account for the singlet (doublet) nature of e_R (l_L) in scattering rates; (ii) the range of temperatures where M_δ (M_H) can be neglected: note that the asymmetry is generated at some scale well above T_{EW} through scattering processes mediated by δ^+ (H).

We also expect that, when both y_ν and y_R take values relevant for leptogenesis, there is more freedom to generate large asymmetries, e.g. because of the presence of extra CP-violating phases, and the requirement of a strong degeneracy between M_2 and M_3 could be relaxed. As in the case of y_ν, one cannot raise too much the value of y_R-entries, as the coherence of sterile neutrino oscillations requires the scattering rate with the SM leptons to be out-of-equilibrium at the time of oscillations; this should translate in a bound similar to the one in Eq. (16). A numerical study is needed to establish more precisely the allowed parameter space, and the differences with respect to the νMSM one. The effect of the two sets of Yukawa couplings could be dramatic, as is the effect of a third GeV scale sterile neutrino [42,44], or of a second Higgs doublet [16].

There is no constraint on y_R coming from the seesaw relation, so one can take couplings smaller than in Eq. (17), slightly enlarging the region of parameters of leptogenesis.

Perhaps more importantly, one can take $(y_R)_{\alpha i}$ significantly larger than the right-hand side of Eq. (17), with no need of fine-tuning to keep m_ν small. Note that the coupling y_R does not induce any mixing with active neutrinos. Indeed, direct searches of active-sterile mixing are presently sensitive to y_ν-entries much larger than in Eq. (17): if the seesaw parameters take their natural values, no direct signal of active-sterile mixing is expected.

Finally, the three-body decay rate of N_i through the coupling y_R is proportional to $|(y_R)_{\alpha i}(y_L)_{\beta\gamma}|^2/M_\delta^4$, to be compared with $|(y_\nu)_{\alpha i}|^2/(M_i^2 v^2)$ for a decay through the mixing with active neutrinos. Since the latter is typically much faster, the nucleosynthesis bound applies to y_ν only, in the form of Eq. (18). Then, leptogenesis through y_R could work even for $M_i < 0.1$ GeV, but in this region nucleosynthesis demands large values for the y_ν-entries, that require strong cancellations in the seesaw.

5 Experimental signatures for N_1 and δ^+

Two important differences between our model and the νMSM are the additional region of the parameter space where the dark matter constraint can be satisfied and the presence of the extra charged scalar δ^+. They both give rise to new experimental signatures that may allow to distinguish one model from the other.

5.1 Dark matter indirect detection

A multi-keV dark matter neutrino is inherently unstable and decays into three light neutrinos at tree-level ($N_1 \to 3\nu_\alpha$) and radiatively into a light neutrino and a photon ($N_1 \to \nu_\alpha\gamma$). This radiative decay produces an X-ray line at $E_\gamma \simeq M_1/2$ that can be searched for and used to constrain the model or to help determine its parameter space [45]. The presence of the additional scalar δ^+ does not affect the decay modes of the dark matter neutrino, which are still determined by its mass and its mixing with the active neutrinos, just as in the νMSM, but it modifies the regions that are consistent with the dark matter constraint, allowing, in particular, for smaller mixing angles. This fact has important implications, as we show next.

Recently, the detection of an unidentified spectral line at about 3.5 keV has been reported from two independent data sets [46,47]. Arguments in favour or against the dark matter decay interpretation of the signal can be found in [48–51]. If confirmed, that signal would provide compelling evidence for keV-scale dark matter and, in particular, for dark matter in the form of sterile neutrinos. Within that framework, the signal can be explained if $M_1 \simeq 7$ keV and $\sin^2 2\theta_1 \simeq 5 \times 10^{-11}$. In the minimal scenario for sterile neutrino dark matter, where they are produced non-resonantly, such parameters lead to a

relic density way below the range determined by cosmological observations. Thus, an additional source of sterile neutrinos is required. A simple possibility for that new source are the decays of the δ^+ particle, as explained in the previous section. These decays allow to decouple the dark matter production (determined by the y_R couplings) from the dark matter decay (determined by y_ν via mixing angles), opening new viable regions consistent with all bounds. Specifically, the freeze-in production of sterile neutrinos that we have examined in this paper enables to explain the tentative 3.5 keV line and to simultaneously account for the observed relic density, even in the absence of a large, primordial lepton asymmetry. If that line signal turns out to be spurious, X-ray observations will continue to be the main way in which the dark matter sector of this model can be tested in the foreseeable future.

5.2 Charged scalar searches at colliders

The singly-charged isosinglet scalar δ^+ could be directly produced at colliders, if its mass were within their energy range. The dominant production channel is the Drell–Yan process $\psi\overline{\psi} \to \gamma/Z \to \delta^+\delta^-$, with $\psi = e$ at LEP and $\psi = q$ at Tevatron and LHC, with a partonic cross-section given e.g. in [52]. The δ^+ decays into one anti-lepton and one anti-neutrino, either through the Yukawa coupling $(y_L)_{\alpha\beta}$, or $(y_R)_{\alpha i}$ when $M_i < M_\delta$.[6]

The standard experimental searches for charged scalars usually assume production and decay modes different from those above. For a singly-charged scalar H^+, the decay channel into lepton plus missing energy has been analyzed, but only for an isodoublet produced from a top-bottom vertex, as in type II two-Higgs doublet models, allowing for an effective background reduction [53,54]. Indeed, the signal over background ratio is significantly smaller in the case of δ^+; an analysis of the cuts required to maximize the signal has been presented in section 5.2 of [55].

Presently, it appears that the best way to constrain M_δ is to use the experimental searches for supersymmetric particles. Since δ^+ and a right-handed slepton have the same gauge quantum numbers, when the latter is directly pair-produced and decays into a lepton and a light neutralino, it behaves very much as the former, and the same bounds apply to both particles. Recently, that setup was studied in the context of simplified supersymmetric models, for the case of a right-handed s-electron degenerate with a right-handed s-muon, decaying with branching ratio 100 % into electron and muon, respectively, plus missing energy [56,57]. Assuming a similar efficiency for electron and muon detection and reducing

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[6] In principle these couplings can all be very small, possibly leading to dominant three-body decays into $l^+l^-W^+$ [23]. In this case δ^+ can be sufficiently long-lived to appear as a curved charged track across the whole detector.

the signal by a factor two, to take into account that the δ^+ production cross-section is the same as for a single right-handed slepton, we can extract a 95 % C.L. lower bound $M_\delta \gtrsim 200$ GeV at ATLAS [56] (see also the auxiliary figures 11a and 12a available online[7]) and $M_\delta \gtrsim 160$ GeV at CMS [57].

Note, however, that these bounds apply as long as the δ^+ branching ratio into taus is negligible. The pair-production of s-taus decaying into taus is presently unconstrained at the LHC (see figure 9 in [58]). As a consequence, if δ^+ decays mostly into taus, its mass can be significantly smaller, say down to the LEP bound ~ 100 GeV. Indeed, in minimal flavour models the δ^+ Yukawa couplings are hierarchical as the SM Yukawa couplings (see Sect. 6), and a large branching ratio into taus is expected.

6 Flavour structure of the model

We accomplished our goal of demonstrating that this model can account for dark matter, neutrino masses, and baryogenesis. That is, we have determined, in the previous sections, that there exist regions in the parameter space of the model where the masses of the new particles and their Yukawa couplings are such that these three issues are simultaneously explained. One may wonder, nonetheless, whether those values of the masses and Yukawa couplings have a generic flavour structure, or if they require a special tuning, indicating that some flavour symmetries are operative. Previous works along this line include [59–63], which focused on the generation of the keV scale in neutrino models, and on the flavour structure of the νMSM. In this section, we do not aim to construct complete flavour models, but rather to recap the order of magnitude of the parameters needed phenomenologically, and to suggest some rationale to explain them.

In view of the hierarchical values of masses and couplings that are required for dark matter and leptogenesis, it is useful to describe the required flavour structure in terms of a $U(1)_F$ family symmetry, with different charges assigned to the various fermions. In this framework, the coefficient of each fermion bilinear $\psi_{Li}\psi_{Lj}$ is suppressed by a power $q_i + q_j$ of a small parameter ϵ (we take the fermion $U(1)_F$-charges to be positive). One may assume some underlying flavour dynamics, that generates $\epsilon = \langle\phi\rangle/\Lambda_F \ll 1$, that is, the vev of a spurion field with $q_\phi = -1$, over the cutoff of the flavour theory.

We need $M_1 \ll M_2 \sim M_3$. Note that in the νMSM a strong degeneracy of M_2 and M_3 is required, while in our scenario this constraint is relaxed. The straightforward way to realize this pattern is to take $q_{N1} > q_{N2} = q_{N3}$, so that $M_1/M_2 \sim \epsilon^{2n}$, with $n \equiv q_{N1} - q_{N2}$. For N-oscillation

(N-decay) leptogenesis, one needs $\epsilon^{2n} \sim 10$ keV/GeV $= 10^{-5}$ ($\epsilon^{2n} \lesssim 10$ keV/TeV $= 10^{-8}$). This charge assignment implies automatically a hierarchy $\epsilon^n : 1 : 1$ among the three columns of the matrix y_R (as well as of y_ν), in other words, $(y_R)^2_{\alpha i}/(y_R)^2_{\alpha j} \sim M_i/M_j$. Since this relation is not satisfied in certain regions of parameters relevant for dark matter, leptogenesis and neutrino masses, an additional flavour symmetry must be operative. The Yukawa couplings can be further suppressed by introducing parity symmetries $Z_2^{(i)}:N_i \rightarrow -N_i$, that allow for M_i but forbid $(y_\nu)_{\alpha i}$ and $(y_R)_{\alpha i}$; these couplings must then be proportional to a small symmetry-breaking parameter ϵ_i. To suppress $(y_\nu)_{\alpha i}$ and not $(y_R)_{\alpha i}$ (or vice versa), one may argue that under the same parity the combination $e_{R\alpha}\delta^+$ (or $l_{L\alpha}H$) is also odd.[8]

Coming to the $U(1)_F$ charges of the three families of lepton doublets and singlets, $q_{L\alpha}$ and $q_{R\alpha}$, first of all they determine the hierarchy of the charged lepton masses, $m_e : m_\mu : m_\tau \sim \epsilon^{q_{Le}+q_{Re}} : \epsilon^{q_{L\mu}+q_{R\mu}} : \epsilon^{q_{L\tau}+q_{R\tau}}$. In turn, the hierarchy among the rows of y_R (y_ν) is determined by the charges $q_{R\alpha}$ ($q_{L\alpha}$). Note that the charges of $l_{L\alpha}$ and $e_{R\alpha}$ are important for charged lepton and active neutrino masses, as well as for leptogenesis from N-oscillations; however they are not very relevant for dark matter production nor for leptogenesis from N-decays.

Let us confront these simple flavour symmetries with the values of the parameters needed for dark matter, neutrino masses and leptogenesis. The freeze-in of the desired amount of N_1 from δ^+ decays requires $\sum_\alpha |(y_R)_{\alpha 1}|^2 \simeq 2.5 \times 10^{-24} M_\delta/M_1$ (see Eq. (10)). The N_1-production from active-sterile oscillations is negligible for $|(y_\nu)_{\alpha 1}| \lesssim 10^{-13}(M_1/\text{keV})$ (barring a large primordial lepton asymmetry). To generate large enough active neutrino masses, Eq. (17) must be satisfied. Clearly N_1 does not contribute significantly to m_ν because of the dark matter constraint, therefore N_2 and N_3 are responsible to generate the atmospheric and solar mass scales. Since $\sqrt{\Delta m^2_{sol}/\Delta m^2_{atm}} \simeq 0.2$, Eq. (17) must hold for both $i = 2$ and 3, up to a factor of a few. Coming to leptogenesis, the two mechanisms discussed in Sect. 4 correspond to two very different regions of parameters that we discuss in turn.

In the case of leptogenesis from N-decays through the y_R-interaction, one requires $M_{2,3} > $ TeV, as well as $|(y_R)_{\alpha 3}| > 10^{-3}\sqrt{M_3/M_2}$ for at least one α, and $|(y_R)_{\alpha 2}| < 10^{-4}\sqrt{M_2/(10^8 \text{ GeV})}$ for all α. For definiteness, let us take $M_1 \sim$ keV, $M_\delta \sim$ TeV and $M_{2,3} \sim 100$ TeV (leptogenesis through y_ν is irrelevant at such low scales, barring resonances). Taking $\epsilon = 0.1$ and the $U(1)_F$ charges $q_{N1} = 7$,

7 https://atlas.web.cern.ch/Atlas/GROUPS/PHYSICS/PAPERS/ SUSY-2013-11/.

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8 The other way around, one could also enhance the hierarchy among the M_i relatively to the hierarchy among the columns of y_ν and y_R, by introducing a (family-dependent) lepton number $U(1)_L$, that is conserved in the Yukawa couplings and broken by two units in the Majorana mass term. We will not need such a symmetry in the following.

$q_{N2} = q_{N3} = 2, q_{R\tau} = 0 \geq q_{R\mu}, q_{Re}, q_{L\tau} = 2 \geq q_{L\mu}, q_{Le}$, we can reproduce the correct size of the M_N, y_R and y_ν entries, except for the second column of y_R, that must be further suppressed by $\epsilon_{R2} \lesssim 10^{-3}$, and the first column of y_ν, to be further suppressed by $\epsilon_{\nu1} \lesssim 10^{-4}$. As discussed above, ϵ_{R2} ($\epsilon_{\nu1}$) can be associated to the breaking of N_2-parity (N_1-parity). Note that here $M_{2,3} \sim \epsilon^4 \Lambda_L$, with a lepton number violation scale $\Lambda_L \sim 10^9$ GeV.

In the case of leptogenesis from N-oscillations coupled to the SM through y_R, one requires for $i = 2, 3$, $M_i < 100$ GeV, $|(y_R)_{\alpha i}| M_i \lesssim 10^{-5}$ GeV, as well as $|(y_R)_{\alpha i}| > 10^{-7}$ for some flavour α and $< 10^{-7}$ for some different flavour β. For definiteness, let us take $M_1 \sim 10$ keV, $M_{2,3} \sim 10$ GeV and $M_\delta \sim$ TeV. Then, one can reproduce the correct size of the M_N and y_R entries taking $\epsilon = 0.1$ and $U(1)_F$ charges $q_{N1} = 8$, $q_{N2} = q_{N3} = 5$, $q_{R\tau} = 1$, $q_{R\mu} = 2$ and $q_{Re} = 3$, and charged lepton masses require $q_{L\tau} = 1 \geq q_{L\mu}, q_{Le}$. Coming to the structure of y_ν, one needs an extra source of suppression for $(y_\nu)_{\alpha 1}$, by a factor $\epsilon_{\nu1} \lesssim 10^{-3}$, to avoid N_1-overproduction from active-sterile mixing; as before, $\epsilon_{\nu1}$ can be understood as the breaking parameter of N_1-parity. We note that, for the chosen values of parameters, both y_R and y_ν couplings to $N_{2,3}$ are relevant for leptogenesis. Here lepton number is broken at the scale $\Lambda_L \sim M_{2,3}/\epsilon^{10} \sim 10^{11}$ GeV.

In summary, the regions of parameters where our scenario is effective require a well-defined hierarchy of masses and couplings, that point to specific but relatively natural flavour structures, with no need of extreme tuning of parameters.

7 Conclusions

We have shown that, when the SM is extended with three sterile neutrinos N_i and a charged scalar δ^+, baryogenesis via leptogenesis and the production of dark matter can both be mediated by δ^+.

The dark matter candidate in this scenario is the lightest sterile neutrino N_1, with a mass M_1 larger than about a few keVs. Such sterile neutrinos are produced by the decays of δ^+ while it is in thermal equilibrium in the early Universe. We showed that this novel mechanism, which is an example of freeze-in dark matter production, can account for the observed dark matter density. Moreover, since the resulting relic density does not depend on the active-sterile mixing angle θ_1, the constraints from X-ray data are alleviated. The recent hint of a signal at 3.5 keV can be easily explained by N_1-decays, for $\theta_1 \sim 3 \times 10^{-6}$. Alternatively, for a vanishingly small θ_1, the dark matter mass M_1 can be raised to the MeV scale, or beyond.

Leptogenesis can be realized through the coupling y_R between δ^+ and $N_{2,3}$ in two different ways, either via N-decays or via N-oscillations. Both mechanisms are well-

known in the context of the SM plus sterile neutrinos only, as they can both proceed through the neutrino Yukawa coupling y_ν. The latter is constrained by light neutrino masses, restricting significantly the allowed parameter space and the testability of leptogenesis. In contrast, in our scenario y_R is not involved in the neutrino mass generation. As a consequence, the N-decay mechanism is successful for $M_{2,3}$ as small as a few TeVs, with no need of a resonant enhancement. The parameter window for N-oscillation leptogenesis is similar when one employs the coupling matrix y_R instead of y_ν: one needs roughly $M_{2,3} \sim$ GeV and y_R-entries smaller than 10^{-6}. However, y_R introduces additional sources of CP-violation, and moreover it is not constrained neither by light neutrino masses nor by nucleosynthesis. Thus one can largely relieve the fine-tuning of parameters necessary in the case with y_ν only.

Neutrino masses are generated by a low scale seesaw mechanism involving the two heavier singlets, $N_{2,3}$, the N_1-contribution being negligible. In this way one can accommodate the current data on neutrino masses and mixing angles, with the prediction of an almost massless lightest neutrino. At colliders such as the LHC, this scenario can be probed and constrained via the direct production of δ^+. This minimal extension of the SM thus provides a simple and testable scenario to explain neutrino masses, the dark matter, and the baryon asymmetry of the Universe.

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