WIMPy baryogenensis with sterile neutrinos

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Abstract. In this talk we propose a mechanism for baryogenensis from particle decays or annihilations that can work at the TeV scale. Some heavy particles annihilate or decay into a heavy sterile neutrino N (with $M \gtrsim 1$ TeV) and a “light” one $\nu$ (with $m \ll 100$ GeV), generating an asymmetry among the two helicity degrees of freedom of $\nu$. This asymmetry is partially transferred to Standard Model leptons via fast Yukawa interactions and reprocessed into a baryon asymmetry by the electroweak sphalerons. We illustrate this mechanism in a WIMPy baryogenensis model where the helicity asymmetry is generated in the annihilation of dark matter. This model connects the baryon asymmetry, dark matter, and neutrino masses.

1. Motivation
The observed dark matter (DM) and baryon energy densities of our Universe are comparable, $\Omega_{DM} \sim 5 \Omega_B$ [1], and both require physics beyond the Standard Model (SM) to be explained. The conventional explanations are unrelated, and often involve very different scales of new physics. The most popular candidates for DM are Weakly Interacting Massive Particles (WIMPs), which provide the so-called “WIMP miracle”: the thermal relic abundance of a stable WIMP is naturally of the order of the observed $\Omega_{DM}$. Regarding the baryon abundance, the Sakharov conditions to generate dynamically the BAU can be fulfilled in a variety of extensions of the SM at very different energy scales, ranging from below the electroweak to the Planck scale.

However the possibility that the BAU is also generated at low temperature, $T \sim$ TeV scale, has received much attention, and it is motivated by several reasons: 1) Mechanisms for baryogenensis at the TeV scale are potentially testable in present or near future experiments, such as LHC, neutrinos oscillations and lepton flavour violating charged lepton decays [2, 3, 4] 2) They avoid hierarchy problems, and, in supersymmetric scenarios, the tension between the high temperatures required by “standard” leptogenesis ($T \gtrsim 10^8$ GeV) and the low reheating temperature ($T \lesssim 10^5 \sim 10^8$ GeV) needed not to overproduce gravitinos [5]. 3) High scale leptogenesis could be disfavored if lepton number violating processes were to be observed at next run of LHC [6, 7]. 4) If baryogenensis takes place at the TeV scale, it could be related to the WIMP DM freeze-out, providing a common origin of dark and baryonic matter, as the similarity of the corresponding energy densities seem to suggest.

Most models proposing a connection between the dark and baryonic matter abundances involve Asymmetric Dark Matter (ADM): the DM we observe today is due to a particle-antiparticle asymmetry in the dark sector which is somehow tied to the one in baryons (for
extensive reviews see [8, 9]). However in the ADM scenario the “WIMP miracle” is lost. This has motivated several attempts to find a mechanism that preserves the natural DM relic density of a WIMP and at the same time relates the dark and baryonic matter abundances [10, 11]. i.e., a baryogenesis scenario via WIMPs. In the next section, we will briefly review two of them: the generation of the BAU in the out-of-equilibrium decay of an unstable WIMP, whose abundance is naturally of the the same order of the DM one [12, 13], and WIMPy Baryogenesis (WB), in which the BAU is generated directly in the annihilation of a stable WIMP [14].

2. Baryogenesis via WIMPs

It is worth to remark that, since baryogenesis via WIMPs takes place at $T \sim \text{TeV}$ scale, it suffers from an intrinsic problem of all scenarios of thermal baryogenesis from particle decays or annihilations at low scales: fast washout of the asymmetries. On general grounds, CP violation requires both, a CP odd and a CP even phase. The former arises from complex couplings, and the later from the absorptive part of one-loop contributions, at leading order. In turn, the absorptive part is proportional to the amplitude of processes that violate baryon or lepton number [15]. As a consequence, such amplitude cannot be very small, for the CP asymmetry to be large enough. Then, if baryogenesis occurs at low temperature, when the expansion rate of the Universe is slow, these processes are typically very fast and washout the asymmetry.

2.1. Baryogenesis in WIMP decay

In this scenario, there are several WIMPs: at least one of them is stable, and makes up the observed DM energy density, but there are also unstable WIMPs, $\chi_B$, which can generate the observed BAU, provided they have CP and $B$ (or $L$)-violating interactions. If $\chi_B$ is long lived, it first undergoes thermal freezeout, and afterwards its late decay, violating $B$ (or $L$) and CP, triggers baryogenesis [13]. The connection between the DM and baryon energy densities arises because the abundance of the decaying $\chi_B$, being a WIMP, is naturally of the same order as the DM relic abundance, while the baryon asymmetry inherits such miracle abundance from the WIMP parent, up to a suppression factor $\epsilon_{\text{CP}}$, the CP asymmetry in $\chi_B$ decay. As a result, one obtains:

$$\frac{\Omega_B}{\Omega_{DM}} = \epsilon_{\text{CP}} \frac{m_p}{M_{\chi_B}} \frac{\Omega_{\chi_B} \to \infty}{\Omega_{DM}},$$

(1)

where $\Omega_{\chi_B} \to \infty$ is the would-be relic abundance of WIMP $\chi_B$ in the limit it is stable. As long as $\chi_B$ survives thermal freezeout via annihilations of weak strength, the resulting $\Omega_B$ is insensitive to the precise value of its lifetime.

In [12], a TeV scale leptogenesis model is presented in which the BAU is generated in the out-of-equilibrium decay of some SM singlet Dirac fermions. In this case, annihilations and inverse decays freeze out at similar temperatures, leading to abundances that are naturally of order that of WIMPs. A pair of quasi-degenerate fermions is required to obtain a large CP asymmetry, $\epsilon_{\text{CP}} \sim \mathcal{O}(1)$, and thus similar relic densities of baryons and WIMPs.

In order to show the above mentioned washout problem of these TeV scale leptogenesis scenarios, let us denote $N_1$ the decaying heavy particle (a Majorana fermion, for definiteness), which has CP and $L$-violating interactions with SM leptons, $\ell$, and an extra scalar, $H_2$. The CP asymmetry in $N_1$ decays is given by,

$$\epsilon_{\text{CP}} = \frac{\gamma(N_1 \to \ell H_2) - \gamma(N_1 \to \bar{\ell} \bar{H}_2)}{\gamma(N_1 \to \ell H_2) + \gamma(N_1 \to \bar{\ell} \bar{H}_2)},$$

(2)

At least another heavy fermion, $N_2$, also coupled to $\ell$ and $H_2$ is needed in order to obtain a non-zero $\epsilon_{\text{CP}}$, due to the interference among the tree-level and one loop diagrams shown in Figure 1. Notice that in the standard type I seesaw leptogenesis, the couplings that appear in
the CP asymmetry are the same Yukawa couplings that determine the active neutrino masses; as a result, a large enough $\epsilon_{CP}$ can be obtained only for $M_1 \gtrsim 10^6 - 10^9$ GeV, if the heavy sterile neutrinos are hierachical, $M_1 \ll M_2$ [16, 17]. However even in the absence of such relation, as explained before the CP asymmetry $\epsilon_{CP}$ is always proportional to the process at the right of the cut in Figure 1, $H_2 \ell \leftrightarrow \bar{H}_2 \bar{\ell}$, which may washout the lepton asymmetry depending on how fast they are compared to the Hubble rate $H$. Since $H \propto T^2/m_P$, with $T$ the temperature, there is an absolute energy scale given by the Planck mass $m_P$: the lower the scale of baryogenesis compared to $m_P$, the harder it becomes to have at the same time a large enough CP asymmetry and small washouts [18].

The importance of these washouts is shown in Figure 2, where we plot the final baryon asymmetry, $Y_B$ as a function of $M_1$, including washout processes mediated by $N_2$ $H_2 \ell \leftrightarrow \bar{H}_2 \bar{\ell}$, and $\ell \ell \leftrightarrow \bar{H}_2 \bar{H}_2$ (solid blue curve) and without including them (dashed red curve). The Yukawa couplings and $M_2$ are chosen so that $\Gamma_1/H(T = M_1)$ and $\epsilon$, remain constant for all $M_1$, where $\Gamma_1$ is the decay width of $N_1$. As can be seen, for $M_1 \lesssim 10^7$ GeV the baryon asymmetry falls exponentially with decreasing $M_1$. Instead, if the washouts are -incorrectly- not included, $Y_B$ only declines linearly and its value for $M_1$ at the TeV scale is wrong by many orders of magnitude.

We believe there are three ways out to overcome this generic problem [?], allowing for baryogenesis at the TeV scale, and therefore at least one of them should be at work if we are to explain the intriguing similarity $\Omega_B \sim \Omega_{DM}$ via WIMPs:

- **Late decay** The rate of the washout processes associated to the CP asymmetry, like $H_2 \ell \leftrightarrow \bar{H}_2 \bar{\ell}$ in Figure 1, is proportional to $(T/M_2)^a$ when $T \lesssim M_2$, with $M_2$ the mass of the mediator and $a = 2 (4)$ for a Majorana (Dirac or scalar) propagator. Increasing the value of $M_2$ relative to $M_1$ (the mass of the decaying particle $N_1$) does not help, of course, to lower the scale of baryogenesis, because the CP asymmetry and washouts would be reduced by the same amount. However the later the $N_1$'s decay, the smaller the washouts will be at the time the baryon (or lepton) asymmetry is generated. Successful baryogenesis at the TeV scale via this late-decay scenario requires two conditions: a small decay width $\Gamma_1 \ll H(T = M_1)$ and a new process to create the $N_1$'s (given that the inverse decays and related scatterings are tiny), e.g. an exotic gauge interaction [19, 20]. Note that the “new” process must decouple before $N_1$ decays, so that the $N_1$'s disappear via the CP-violating interaction. In this way it is possible to have leptogenesis at $T \sim$ few TeVs (the exact value
Figure 2. Baryon asymmetry, $Y_B$ as a function of $M_1$, including washout processes $H_2 \ell \leftrightarrow \bar{H}_2 \bar{\ell}$, and $\ell \ell \leftrightarrow H_2 H_2$ (solid blue curve) and without including them (dashed red curve). The Yukawa couplings and $M_2$ are chosen so that $\Gamma_1/H(T = M_1)$ and $\epsilon$ remain constant for all $M_1$.

depending on the freeze-out temperature of the sphalerons) and baryogenesis at much lower scales if $B$ is violated perturbatively.

The late decay mechanism was incorporated in [13], where the scale of baryogenesis via the decay of the WIMP $\chi_B$ is only constrained to be above that of Big Bang nucleosynthesis.

- **Resonant baryogenesis**: When there is a pair of almost degenerate particles, the CP asymmetry in decays can be enhanced up to $O(1)$ values [21, 22]. For instance, in the case of two sterile neutrinos $N_i \{i = 1, 2\}$, with masses $M_i$, Yukawa couplings $\lambda_{\alpha i}$ with the SM lepton doublet $\ell_\alpha$, and decay widths $\Gamma_i$, if the resonant condition $\Delta M \equiv M_2 - M_1 = \Gamma_2/2$ holds, the CP asymmetry in $N_1$ decays, $\epsilon_1$, can take values as large as 1/2 independently of the size of the Yukawa couplings of $N_2$, i.e. of $(\lambda^\dagger \lambda)_{22}$. Moreover, for $\Gamma_{1,2} \ll \Delta M \ll M_1$, $\epsilon_1 \propto (\lambda^\dagger \lambda)_{22}/\delta$, with $\delta \equiv \Delta M/M_1$. Therefore it is possible to reduce the “dangerous” washouts taking $(\lambda^\dagger \lambda)_{22}$ small enough, while keeping $\epsilon_1$ sizable by choosing a tiny value for $\delta$. This -so called resonant leptogenesis- mechanism has been widely studied (see e.g. [23]) and it can lead to successful baryogenesis at the TeV scale.

Resonant leptogenesis is required in [12], to obtain a large enough CP asymmetry while suppressing the unwanted washout scattering processes.

- **Massive decay or annihilation products**: Although the CP asymmetry is proportional to the amplitude of processes that violate $B$ (or $L$), the rate of the washouts is not only determined by the amplitude of the corresponding process, but also by the number density of the particles involved in the interaction. If at least one of these particles is massive and becomes non-relativistic during the baryogenesis epoch, the corresponding washout rate will be Boltzmann suppressed. This mechanism was shown to allow for TeV scale baryogenesis from DM annihilation first in [14] and from heavy particle decays in [18].
2.2. Baryogenesis from DM annihilation: WIMPy baryogenesis

The Sakharov conditions for baryogenesis can also be satisfied during the annihilation freezeout of a WIMP-like DM particle. In this alternative scenario, so-called WIMPy baryogenesis (WB), the BAU is generated directly in the annihilation of a stable WIMP [14], so the WIMP miracle is preserved and the observed relation $\Omega_{DM} \sim 5 \Omega_B$ can be obtained with a moderate adjustment of parameters. The phenomenology of several WB models has been studied in [24] and the conditions for generating the observed $\Omega_B$ and $\Omega_{DM}$ via this mechanism analyzed in detail in [25].

In WB the DM, $\chi$, is a weakly interacting massive particle whose relic density is determined by the freeze out of some annihilation process $\chi\chi \rightarrow \Psi f$, with $f$ a SM fermion and $\Psi$ a heavy exotic particle. The amplitude for the process $\chi\chi \rightarrow \Psi f$ contains a CP odd phase coming from complex couplings and a CP even phase from the absorptive part of one loop contributions, therefore it violates CP. Moreover, depending on whether $f$ is a SM quark or a lepton, the interaction $\chi\chi \rightarrow \Psi f$ violates SM B or L, respectively. In this way all Sakharov conditions are satisfied and some baryon or lepton asymmetry is produced in the annihilation of DM.

The minimal content for WIMPy baryogenesis consists of a fermonic DM field, $\chi$, coupled to at least two scalars $S_1, S_2$ (in order to have a non zero CP asymmetry) and a heavy fermion $\Psi$. The asymmetry generated in $\chi\chi$ annihilation is defined as:

$$\epsilon_{CP} = \frac{\gamma(\chi\chi \rightarrow L_i \Psi_i) - \gamma(\chi\chi \rightarrow \bar{L}_i \bar{\Psi}_i)}{\gamma(\chi\chi \rightarrow L_i \Psi_i) + \gamma(\chi\chi \rightarrow \bar{L}_i \bar{\Psi}_i)}. \quad (3)$$

In Figure 3 are depicted the relevant diagrams contributing to the CP asymmetry when the DM annihilates into leptons mainly through the lightest scalar, $S_1$. Successful WIMPy leptogenesis, i.e., correct BAU and WIMP relic density, is obtained for $M_\chi \sim (1 - 10) \text{ TeV}$, $M_\chi < M_{S_1, S_2}$, $0.5 M_\chi \lesssim M_\Psi < 2 M_\chi$ and Yukawa couplings of $O(1)$ [14].

The problem of fast washout of the asymmetry mentioned before is also present in this framework, but it can be overcome by the last mechanism described in the previous section: if $\Psi$ is heavy enough, $m_\Psi \gtrsim m_\chi$ [14], the processes that can potentially washout the asymmetry -mainly $\bar{\Psi} f \leftrightarrow \Psi f$- are Boltzmann suppressed, hence a significant amount of matter asymmetry may survive.

A crucial point for this mechanism to work is the following: the annihilation $\chi\chi \rightarrow \bar{\Psi} f$ also generates an asymmetry in the $\Psi$ sector and it is not trivial to avoid a cancellation of the total matter asymmetry after $\Psi$ disappears from the thermal bath. As a consequence, more complicated models are needed, for example in [14] the $\Psi$ decay into a light hidden sector, while decays into SM particles are forbidden by a $Z_4$ symmetry.

From the Boltzmann equation

In the next section we present a WIMPy leptogenesis model which solves this problems in a simple way, and moreover brings a connection to light neutrino masses.
3. A WIMPy baryogenesis model with sterile neutrinos

3.1. The mechanism

In [26] we propose a WIMPy leptogenesis scenario where the role of $\Psi$ is played by heavy Majorana fermions (subsequently called $N_i$) which lack a conserved charge to store asymmetry, avoiding the complications just mentioned. Therefore no additional fields beyond those participating in the annihilation of DM are needed. It can be realized in both, baryogenesis from DM freeze out and from heavy particle out-of-equilibrium decay. We illustrate this mechanism in a WB model where the DM annihilates into sterile neutrinos, which in turn are responsible for neutrino masses via the type I seesaw. In this way we address the possibility of relating the DM and BAU problems with yet another puzzle that requires physics beyond the SM: neutrino masses.

The basic requirement for generating the BAU is the existence of heavy sterile neutrinos $N (M \gtrsim 0.5 \text{ TeV})$ and also lighter ones $\nu$ (with $m \ll 100 \text{ GeV}$), both of them interacting with the DM via SM singlet scalars. In a first step, an asymmetry among the two helicity degrees of freedom of $\nu$ is generated via CP violating WIMP annihilations into sterile neutrinos, $\chi\chi \rightarrow N\nu$. Then, this asymmetry is transferred to SM leptons via fast Yukawa interactions, which should be in equilibrium prior to the electroweak phase transition to ensure that a baryon asymmetry is also induced by the sphaleron $(B+L)$-violating interactions. Moreover, since no asymmetry accumulates in the heavy sector $N$, some of the requisites of the original WB models- a $Z_2$ symmetry and a light sterile dark sector- are automatically avoided. Note that in the limit $m \rightarrow 0$ there is a global $U(1)$ symmetry. This implies the existence of a -perturbatively-conserved lepton number which coincides with the helicity for the $\nu$'s, i.e. a $\nu$ with positive (negative) helicity has lepton number $L_\nu = 1 (-1)$.

Given that the mass scale of the sterile neutrinos is unconstrained and its origin unknown, it is worthwhile to explore the consequences of having several mass scales in the sterile sector without any theoretical prejudice. Therefore we first adopt a purely phenomenological perspective, illustrating the proposed mechanism by means of a minimal model.

In the minimal scenario, the SM is extended with some singlet real scalars, $S_a$, and Majorana fermions, $\chi, N_i, \nu_j$, with $\{a, i, j = 1, \ldots\}$, together with a discrete $Z_2$ symmetry to ensure the stability of the DM. The DM candidate $\chi$ is the only odd particle under $Z_2$. It could also be a Dirac singlet, but we choose it Majorana to minimize the number of new degrees of freedom. The $N_i$ and $\nu_j$ are sterile neutrinos with “high” ($M_i \sim \mathcal{O}(\text{TeV})$) and “low” ($m_j \ll 100 \text{ GeV}$) masses, respectively. In the basis which yields a diagonal Majorana mass matrix with real and positive entries, the most general renormalizable Lagrangian with the given fields and symmetries reads

$$-\mathcal{L} = -\mathcal{L}_{\text{SM}} - \mathcal{L}_{\text{kin}} + V(S_a, H) + \frac{1}{2} \left\{ m_\chi \bar{\chi}\chi + M_i \bar{N}_i N_i + m_{ij} \bar{\nu}_i \nu_j \right\} + \frac{1}{2} \left\{ \lambda_{\nu a} S_a \bar{\nu}_i P_R \nu_j + \lambda_{N a i j} S_a \bar{N}_i P_R N_j + \lambda_{\nu a i j} S_a \bar{\nu}_i P_R \nu_j + h.c. \right\} + \left\{ \lambda_{\nu a i j} S_a \bar{N}_i P_R \nu_j + h_{\nu a i j} \bar{H} \ell_\alpha P_R N_i + h_{\nu a i j} \bar{H} \ell_\alpha P_R \nu_j + h.c. \right\}$$

(4)

where there is an implicit sum over repeated family indices, $\ell_\alpha$ are the leptonic $SU(2)$ doublets, $H$ is the Higgs field ($\bar{H}_2 = i\tau_2 H_2^\ast$, with $\tau_2$ Pauli’s second matrix), $V(S_a, H)$ is the scalar potential and $P_{R,L} = (\pm \gamma_5)/2$ are the chirality projectors. Latin indices denote sterile neutrinos while Greek indices refer to the SM lepton doublets. All Majorana fields $\xi (\xi = \chi, N_i, \nu_j)$ satisfy $\xi^\dagger = \lambda_\xi \xi$, with $\lambda_\xi$ a phase factor. Notice that the Yukawa matrices $\lambda_{N a i}, \lambda_{\nu a}$ are symmetric.

1 Although in this model the species $\nu_j$ and $N_i$ differ only in their masses, we denote them by different symbols to emphasize their distinct roles for leptogenesis.
The key for having baryogenesis in this model is that $m_j \ll T_{sfo}$ at least for one species $\nu_j$, where $T_{sfo} = \mathcal{O}(100 \text{ GeV})$ is the temperature at which the electroweak sphalerons freeze out. Then an asymmetry among the two helicity degrees of freedom of $\nu_j$, created from $\chi\chi$ from the annihilation of DM $\nu_\chi \nu_\chi$ couplings of $\nu_j$, $h_{\nu\nu\nu}$, are large enough, the helicity asymmetry in the $\nu_j$ is efficiently transferred to the SM lepton sector. In turn, this is partially transformed into a baryon asymmetry by the sphaleron processes. Once these decouple at $T_{sfo}$, the BAU is frozen. Notice that if $m_j \neq 0$, the helicity depends on the reference frame. We will be always working in the thermal bath rest frame.

This baryogenesis scenario requires at least one species of $N'$s and $\nu'$s, and two real scalars to have CP violation, $S_1$ and $S_2$ (actually we will explain later that there can be a CP odd phase with just one scalar, contrary to previous WB models, but the amount of CP violation is most likely too small to have successful baryogenesis). Next we indicate the approximate range to have CP violation, $S_{\text{frame}}$. We will be always working in the thermal bath rest frame.

The allowed region in the parameter space of $m_{\chi}, m_{S_a}, M_i$, $\lambda_{\chi a}$ and $\lambda_{\nu a}$ is very similar to previous models of WB [14, 24, 25] (with $N_i$ playing the role of the heavy exotic annihilation product): very roughly it consists of masses above $\sim 1$ TeV and $\mathcal{O}(1)$ couplings. On the other hand, if helicogenesis occurs in the decay of the lightest scalar, $S_1$, the conditions for generating

- $m_\chi$: To generate the asymmetry before sphalerons freeze out, the DM has to start annihilating well above $T_{sfo}$, hence $m_\chi \gtrsim 1$ TeV.
- $m_{S_a}$: Although the asymmetry could also be produced in the decays of $S_a$ (see [14]), we are interested in the case that the asymmetry is mainly produced in the annihilation of DM, hence we impose that the masses of the singlet scalars, $m_{S_a}$, are $m_{S_a} \gtrsim m_\chi$ (in this way the CP conserving annihilation channel $\chi\chi \rightarrow S_aS_a$ is negligible).
- $M_i$: The heavy $N_i$ are introduced to have a Boltzmann suppression $e^{-M_i/T}$ of washouts that can be very fast when baryogenesis occurs at low temperatures (see [14] and [25, 18] for detailed discussions on this point). For this Boltzmann suppression to be efficient $M_i \gtrsim (0.5 - 1) m_\chi$. In addition, $M_i < 2m_\chi$ to allow for DM annihilations when $\chi$ becomes non-relativistic.
- $m_j$: To create an helicity asymmetry in the $\nu$-sector it is necessary that $m_j \ll T_{sfo}$.
- $\lambda_{\chi a}$, $\lambda_{\nu a}$: These couplings must be $\mathcal{O}(1)$ for having enough CP violation and a correct DM relic abundance. More precisely, it is the imaginary part of $\lambda_{\chi a}$ that has to be large, so that there is a sufficiently fast, not velocity-suppressed annihilation rate.
- $\lambda_{\nu aij}$: They induce washouts of the helicity asymmetry that are not Boltzmann-suppressed, e.g. via the reaction $\nu_j^+\nu_j^- \leftrightarrow \nu_j^-\nu_j^-$. For these processes to be slow enough $|\lambda_{\nu aij}m_\chi/m_{S_a}| \ll 10^{-3}$.
- $\lambda_{S aij}$: They do not play an important role because the corresponding processes are Boltzmann suppressed, hence they are unconstrained.
- $h_{\nu aij}$: It is crucial that there be at least one fast Yukawa interaction between the $\nu_j$ and $\ell_\alpha$, and therefore at least one coupling $h_{\nu aij} \gtrsim 2 \times 10^{-7}$ [27, 28].
- $h_{N aij}$: They mediate washout processes like $\ell_\alpha H \leftrightarrow \ell_\alpha H$, which should be slow at $T \gtrsim T_{sfo}$. As for the $\lambda_{\nu aij}$ couplings, this requirement is satisfied for $|h_{N aij}m_\chi/M_i| \ll 10^{-3}$.

In addition the heavy singlet sector must be populated at $T \gtrsim m_\chi$. This can be achieved by some fast interaction connecting the sterile and SM sectors, like one among the $S_a$ and the Higgs or the Yukawa interactions between the $N_i$ and $\ell_\alpha$.

The allowed region in the parameter space of $m_\chi, m_{S_a}, M_i, \lambda_{\chi a}$ and $\lambda_{\nu a}$ is very similar to previous models of WB [14, 24, 25] (with $N_i$ playing the role of the heavy exotic annihilation product): very roughly it consists of masses above $\sim 1$ TeV and $\mathcal{O}(1)$ couplings. On the other hand, if helicogenesis occurs in the decay of the lightest scalar, $S_1$, the conditions for generating
the $\nu$ helicity asymmetry are analogous to those for standard leptogenesis at the TeV scale, basically $\lambda_{\alpha ij} \lesssim 10^{-7}$, $m_{S\alpha} \gtrsim 1$ TeV, and $M_t \gtrsim 0.5$ TeV [18]. Hence we are not going to develop these issues further.

The connection of leptogenesis via helicitogenesis with light neutrino masses imposes new constraints:

(i) The $N$’s and $\nu$’s must decay before Big Bang Nucleosynthesis (BBN) to avoid observational constraints. This requirement is not difficult to accomplish for the heavy neutrinos $N_i$. One possibility is that $h_{N\alpha i}$ be non-negligible to allow for the decay $N_i \rightarrow \ell_\alpha H$, but at the same time small enough for the washout processes like $\ell_\alpha H \leftrightarrow \ell_\alpha H$ to be slow at $T \gtrsim T_{sfo}$. This condition is easy to satisfy given that the rate of this last process is $\propto h_{N\alpha i}^2$, while the rate of the former is $\propto h_{N\alpha i}^2$. Another possibility could be to choose $\lambda_{\nu\alpha ij}$ large enough to induce three body decays like $N_i \rightarrow \nu_j \nu_j \nu_j$, but not as large as to have fast washouts $\nu_j^+ \nu_j^- \leftrightarrow \nu_j^- \nu_j^-$ (again note that the rate of this process is $\propto \lambda_{\nu\alpha ij}^2$ while the rate of $N_i \rightarrow \nu_j \nu_j \nu_j$ is $\propto \lambda_{\nu\alpha ij}^2$).

The main decay modes of the (lightest) $\nu_j$ are $\nu_j \rightarrow \nu_\alpha \nu_\beta \nu_\beta$, $\nu_j \rightarrow \nu_\alpha e^- \bar{e}^+$, $\nu_j \rightarrow \nu_\alpha q_\beta \bar{q}_\beta$, via $Z$ exchange, $\nu_j \rightarrow e^- e^+ \nu_\beta$, $\nu_j \rightarrow e^- q_\beta \bar{q}_\beta$ via $W$ exchange, and the corresponding CP-conjugate processes, where $e_\alpha$ denotes the charged leptons $e, \mu, \tau$ and $q_\beta$ stands for the SM quarks, except the top. The $\nu_j$ decay width is given by

$$\Gamma_j = \frac{G_F^2 m_j^3}{192\pi^3} \sum_{\alpha,\beta} A_{\alpha\beta} |h_{\nu\alpha j} v|^2,$$

where the sum extends over the kinematically allowed decay channels, $v = \langle H \rangle = 174$ GeV is the Higgs vev, $G_F$ is the Fermi constant, and $A_{\alpha\beta}$ are $O(1)$ coefficients that depend on the number of degrees of freedom associated to each mode. Using the above equation, with at least one $h_{\nu\alpha j} \sim 10^{-7}$ we find that $m_j \gtrsim 1$ GeV, in order for $\nu_j$ to decay before BBN. Moreover, even if the decay of the $\nu_j$ occurs safely before BBN, it may lead to an increase of entropy density after the electroweak phase transition, which would dilute the baryon asymmetry. We have checked that this entropy increase is negligible for $m_j \gtrsim 10$ GeV, when the $\nu_j$ decays at $T \sim 500$ MeV, before the QCD phase transition. However it can be a concern for lower masses, $m_j \sim 1$ GeV. In this case, the decay occurs after the QCD phase transition and the increase in entropy density can be up to order 10. This implies that the baryon asymmetry originally produced should be an order of magnitude larger, which could require Yukawa couplings close to the perturbative limit.

(ii) Constraints from neutrino masses:

The seesaw contribution of $\nu_j$ to the light neutrino masses is given by $(m_L)_{\alpha\beta} \sim h_{\nu\alpha j} h_{\nu\beta j} \frac{v^2}{m_j}$. Taking into account that $h_{\nu\alpha j} \gtrsim 2 \times 10^{-7}$, one has that $(m_L)_{\alpha\beta}[eV] \gtrsim 1/m_j[GeV]$. The strongest constraints on the absolute scale of neutrino masses are derived from cosmological observations, via their contribution to the energy density of the Universe and the growth of structure [1]. Since these bounds are very sensitive to the assumptions about the expansion history of the Universe and to the data included in the analysis, we choose the conservative upper bound on the sum of light neutrinos masses of roughly 1 eV, obtained by combining CMB and large scale structure data when including several departures from the $\Lambda$CDM model [29]. This bound implies that $m_j \gtrsim 3$ GeV, unless there is a fine tuning among the phases of the Yukawa couplings of different species of $\nu$’s, so that they give big contributions to $(m_L)_{\alpha\beta}$ with opposite signs that cancel each other. The atmospheric mass scale, 0.05 eV, can be naturally obtained with $m_j \sim 20$ GeV, thus the two conditions for our mechanism to work, $m_j \ll T_{sfo} \sim 100$ GeV and $h_{\nu\alpha j} \gtrsim 2 \times 10^{-7}$ are compatible with the observed
light neutrino masses. Notice that these two requirements are also needed when the helicity asymmetry in the singlet neutrinos is generated via neutrino oscillations [30].

Analogously, the heavy singlets $N_i$ also contribute to light neutrino masses an amount $h_{N\alpha i}h_{N\beta i}v^2/M_i$. Barring accidental cancellations, $h_N \lesssim (10^{-5} - 10^{-6})\sqrt{M_i}/T_{\text{TeV}}$ is consistent with present data, and also allows for the $N_i$ decay before BBN.

One may worry that the separation of the mass scales $m_j \ll M_i$ is not stable under radiative corrections, since there is not any symmetry protecting the small masses. In fact, $\nu_j$ self-energy diagrams with virtual $S_a$ and $N_i$ will induce Majorana masses for the $\nu_j$ at one loop of order

$$m^{1\text{-loop}} \sim \frac{(\lambda_{aij})^2}{16\pi^2} M_i \log \left( \frac{M_i^2}{m_{S_a}} \right).$$

Thus for $\lambda_{aij}$ of $O(1)$, generically required for WB, we expect $m^{1\text{-loop}} \sim 10^{-2} M_i$, which does not upset the condition $m \ll 100 \text{ GeV}$ for $M_i$ of order few TeV. Indeed, this loop contribution is naturally of the correct size for helicitogenesis to work.

Notice that since $S_a$ are real scalar singlets, the potential $V(S_a, H)$ in eq. (4) will contain trilinear terms of the form $\mu_a S_a H^\dagger H$. After electroweak symmetry breaking such terms induce vevs $\langle S_a \rangle \neq 0$, which in turn generate a mixing among the light, $\nu$, and heavy, $N$, neutrino species. We assume that the $\mu_a$ parameters are not too large, therefore the corrections to the $\nu$ masses still satisfy the requirement $\Delta m \sim (\lambda_{aij}(S_a))^2/M_i \ll T_{sfo}$ and the mixing does not affect the proposed baryogenesis mechanism.

3.2. Evolution equations for the helicity asymmetry

Next we calculate the helicity asymmetry in the $\nu_j$ sector, and its partial transformation into a baryon asymmetry. For simplicity we will consider only one species of $N$'s and $\nu$'s, $N \equiv N_1$ and $\nu \equiv \nu_1$, and hence we will omit the indices associated with the $\nu$ and $N$ sectors. As noted in [31], in the thermal bath rest frame isotropy implies that the spin density matrix is diagonal in the helicity basis. This allows to write a set of BEs for the populations of $\nu^+$ and $\nu^-$ involving no coherences. Actually, the quantity of interest is the helicity asymmetry $Y_{\Delta\nu} \equiv Y_{\nu^+} - Y_{\nu^-}$, where for any particle $X$ we define $Y_X \equiv n_X/s$ as the number density of $X$ normalized to the entropy density.

The asymmetry $Y_{\Delta\nu}$ originates from interactions in the singlet sector of the model. In turn, this asymmetry is partially transferred to the lepton sector via the Yukawa interactions among $\nu$ and the SM lepton doublets. Finally the electroweak sphalerons transform part of the lepton asymmetry into a baryon one. A fairly good approximation is to consider that these different stages do not happen simultaneously: first $Y_{\Delta\nu}$ is generated while the DM annihilations freeze out and only then the Yukawa interactions and sphalerons act to get the final BAU. In other words, we neglect spectator processes during the generation of the helicity asymmetry. From the results of [32] we expect that this type of approximation is accurate within factors not larger than $\sim 2$. We will also assume that thanks to the $N_i$-interactions described in the previous section, $Y_{N_i}$ follows an equilibrium distribution while the DM is annihilating. Then $Y_{\Delta\nu}$ can be obtained from the following set of BEs (the details of the derivation of these BEs are very

\footnote{The population of $\nu$'s is in kinetic equilibrium due to different fast processes like scatterings with the $N$'s and Yukawa interactions with SM leptons. This allows for an easy integration of the momentum degrees of freedom, leading to simple Boltzmann equations for the number densities.}
similar to those described in, e.g., the appendix B of [25]):

\[
\sigma_{\mathrm{ann}}(\chi\chi \rightarrow \nu N) = \frac{m_{\chi}^4}{64 \pi^4 z} \int_{x_{\mathrm{min}}}^{\infty} dx \sqrt{x} \sigma_R(x m_{\chi}^2) K_1(z \sqrt{x}) ,
\]

where \(x = s/m_{\chi}^2, x_{\mathrm{min}} = \max \left\{ \left( \frac{m_a + m_b}{m_{\chi}} \right)^2, \left( \frac{m_a - m_b}{m_{\chi}} \right)^2 \right\} \), and here \(s\) is the center of mass energy squared. The reduced cross section \(\sigma_R\) is related to the total cross section \(\sigma\) via

\[
\sigma_R(s) = \frac{2 \lambda(s, m_a^2, m_b^2)}{s} \sigma(s) \quad \text{and} \quad \lambda(s, m_a^2, m_b^2) \equiv (s - (m_a + m_b)^2) (s - (m_a - m_b)^2) .
\]

In the Eq. (7) we have neglected the CP asymmetry in the decay of \(S_a\), which is a good approximation if \(m_{S_a} \gtrsim 2m_{\chi}\) because that asymmetry would be washed out very efficiently. Hence the helicity asymmetry is generated mainly in the annihilation of DM and the CP asymmetry per annihilation appearing in Eq. (7), \(\epsilon\), is defined as

\[
\epsilon \equiv \frac{\Delta \gamma (\chi\chi \rightarrow \nu N)}{\gamma (\chi\chi \rightarrow \nu N)} ,
\]

where \(\Delta \gamma (\chi\chi \rightarrow \nu N) = \gamma (\chi\chi \rightarrow \nu^+ N) - \gamma (\chi\chi \rightarrow \nu^- N)\) and \(\gamma (\chi\chi \rightarrow \nu N)\) is the total annihilation rate, \(\gamma (\chi\chi \rightarrow \nu N) = \gamma (\chi\chi \rightarrow \nu^+ N) + \gamma (\chi\chi \rightarrow \nu^- N)\). As long as there are at least two species of scalars, \(S_1\) and \(S_2\), there is a contribution to the CP asymmetry at zeroth order in \(m/m_{S_0}\) (where \(m \equiv m_1\)). Up to \(O(1)\) numerical factors, the reduced cross sections relevant for the calculation of the CP asymmetry, in the limit \(m \rightarrow 0\), are given by [25]:

\[
\Delta \sigma_R (\chi\chi \rightarrow \nu N) = \frac{1}{8\pi^2 s^{3/2}} \left( \frac{m_{\chi}^2}{s - m_{S_1}^2} \right) \times \\
\left\{ 2\lambda_1 \lambda_2 \mathrm{Im} (\lambda_1 \lambda_2) \left| \chi_1^2 \right|^2 \frac{f_S(m_{S_1}) + f_V(m_{S_1})}{s - m_{S_1}^2} - \left| \lambda_2 \right|^2 \frac{f_S(m_{S_2}) + f_V(m_{S_2})}{s - m_{S_2}^2} \right\} \\
- \left| \lambda_1 \lambda_2^2 \right| \left( \frac{2\lambda_1^2 f_S(m_{S_1}) + f_V(m_{S_1})}{s - m_{S_1}^2} - \frac{2\lambda_2^2 f_S(m_{S_2}) + f_V(m_{S_2})}{s - m_{S_2}^2} \right) ,
\]

\(\lambda_1, \lambda_2\) are the eigenvalues of the matrix \(\chi_{\chi\nu}\), and \(f_S(m_{S_i}) \approx f_V(m_{S_i}) \approx f_{\chi}(m_{S_i})\) are the vector and axial-vector couplings. Note that the \(N\)'s are non-relativistic in the relevant epoch for baryogenesis, hence the chiral operators \(P_R N_1\) and \(P_L N_2\) in Eq. (4) can create and destroy any helicity state of \(N\). Therefore the helicity of the heavy sterile neutrinos does not play a major role and the rates are defined summing over the spin degree of freedom of \(N\).

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and

\[ \sigma_R (\chi \chi \rightarrow \nu N) = \frac{1}{8\pi} (s - M^2)^2 \sqrt{1 - 4m_\chi^2/s} \left| \frac{\lambda_{\chi 1} \lambda_1}{s - m_{\chi 1}^2} + \frac{\lambda_{\chi 2} \lambda_2}{s - m_{\chi 2}^2} \right|^2, \]

where \( \lambda_a \equiv \lambda_{a11} \). The loop functions, \( f_S(m_{S_a}) \) and \( f_V(m_{S_a}) \) (that also depend on \( M \equiv M_1 \)), as well as the remaining density rates in Eq. (7), can be found in appendix A of [25].

Notice that the Majorana nature of \( \nu \) and \( N \) imply that the case of decays this means that the particle decaying and the one in the loop can be identical. Therefore this brings a qualitative difference with typical baryogenesis scenarios. In particular, for interesting enough, this type of contributions requires just one species of real scalars, \( S \), and therefore this brings a qualitative difference with typical baryogenesis scenarios. Actually those contributions are suppressed by \( \frac{m}{m} \) (for annihilations) or \( \frac{m}{m_S} \) (for decays), which is a fair small factor given that \( m \ll T_{fo} = \mathcal{O}(100) \) GeV and \( m_\chi, m_S \gtrsim 1 \) TeV. Whether or not it is possible to have successful baryogenesis with this CP asymmetry is an interesting question, although it would probably require handling very large \( \lambda_1 \) couplings.

The second step in the approximation mentioned above is to analyze how the final helicity asymmetry obtained from the BEs (7) and (8), \( Y_{\Delta \nu}^f \), is transferred to the SM lepton sector. This occurs through different fast processes, all of them involving the Yukawa interactions among \( \nu \) and the SM leptons. When \( m = 0 \) it is possible to define a lepton number \( L \) which is conserved by all these reactions, namely \( L = L_{SM} + L_\nu \), where \( L_{SM} \) is the usual left-handed number for SM fields, while \( \nu^+ (\nu^-) \) is assigned \( L_\nu = 1 (-1) \). Then the chemical equilibrium condition for the Yukawa interactions yields \( \mu_\nu - \mu_\ell = \mu_H \), where \( \mu_X \) is the chemical potential of the particle \( X \). Taking into account also the whole set of relations among chemical potentials due to all the fast SM processes (including the electroweak sphalerons), the conservation laws, and the relation among chemical potentials and density asymmetries [33, 32], one gets that \( Y_{\nu}^B = aY_{\Delta \nu}^\nu \). Here \( a \) is a numerical factor whose value lies between \( \sim 1/4 \) and \( \sim 1/3 \) depending on how many independent fast Yukawa interactions there are among the \( \nu_j \) and \( \ell_\alpha \).

Instead, when \( m \neq 0 \), \( L \) is not conserved. However if \( m \ll T \) the rate of processes violating \( L \) is suppressed by \( (m/T)^2 \) with respect to the rate of reactions conserving \( L \). Hence a net lepton asymmetry can be transferred to the SM sector and baryogenesis can be successful as long as \( m \ll T_{fo} \).

### 3.3. \( U(1)_L \) symmetric scenario

A drawback of the minimal model we have described is that there is no justification for the necessary hierarchical spectrum of sterile neutrino masses and couplings. Actually those hierarchies would suit very well if the singlet fields were charged under a conserved lepton number. The assignment \( L = 0, 1, 1 \), and 1/2 to \( N_i, \nu_R, S_a \), and \( \chi \), respectively, would imply that \( m_j = \lambda_{naij} = \lambda_{Naij} = h_{Nai} = 0 \), which perfectly fulfills the requirements of the helicogenesis mechanism. The most general \( L \)-conserving Lagrangian can be written as

\[ -\mathcal{L} = -\mathcal{L}_{SM} - \mathcal{L}_{kin} + V(S_a, H) + m_\chi \chi \chi + \frac{1}{2} M_i \overline{N}_i N_i + \frac{1}{2} \left\{ \lambda_{\chi aR} S_a^\dagger \overline{\chi} \overline{\chi} P_R \chi + \lambda_{\chi aL} S_a^\dagger \overline{\chi} \overline{\chi} P_L \chi + h.c. \right\} + \left\{ \lambda_{aij} S_a^\dagger \overline{N}_i P_R \nu_j + h_{aij} \overline{H}_i P_R \nu_j + h.c. \right\} \]

with

\[ V(S_a, H) = -m_{S_a}^2 S_a^\dagger S_a + \lambda_{Hab} (H^\dagger H) (S_a^\dagger S_b) + \lambda_{abcd} (S_a^\dagger S_b) (S_c^\dagger S_d) + h.c. \quad (15) \]

In this case the \( N_i \) can decay before BBN thanks to the large couplings \( \lambda_{naij} \) and the \( S_a - H \) mixing after electroweak and \( U(1)_L \) symmetry breaking.
The DM field \( \chi \) is now a Dirac fermion, and \( \chi^c = C \chi^T \). Recall that WB requires at least two scalar fields \( S_a \). If \( m_{S_a}^2 > 0 \) the complex scalars acquire a non-zero vev, lepton number gets spontaneously broken and \( \chi \) splits into two Majorana fermions \( \chi_1, \chi_2 \) with masses

\[
m_{\chi_1, \chi_2} = \frac{1}{2} \left\{ \mu_L + \mu_R \pm \sqrt{(\mu_L - \mu_R)^2 + 4m_{\chi}^2} \right\},
\]

where \( \mu_{L,R} \equiv \sum_a \lambda_a \phi_L \phi_R u_a \) and \( u_a = \langle S_a \rangle \). Notice that in this model we do not need an additional \( Z_2 \) symmetry, as it is usually the case to avoid the DM decay: the lightest \( \chi \) is stable because of a \( Z_2 \) symmetry which is an unbroken remnant of the global \( U(1)_L \), as in [34].

The light neutrino masses would be obtained via a double seesaw [35] mechanism; once the electroweak symmetry is also broken, the SM doublet neutrinos \( \nu_\alpha \) and the sterile ones, \( \nu_j, N_i \), mix and the mass matrix in the \( (\nu_\alpha, \nu_j, N_i) \) basis becomes:

\[
\mathcal{M} = \begin{pmatrix}
0 & h_{\nu \alpha} & 0 \\
h_{\nu \alpha}^T v & 0 & \lambda_{\alpha}^T u_a \\
0 & \lambda_a u_a & M
\end{pmatrix},
\]

where \( v = \langle H \rangle \), the matrix elements of \( \lambda_a \) (\( h_{\nu \alpha} \)) are the Yukawa couplings \( \lambda_{\alpha ij} \) (\( h_{\nu \alpha ij} \)), and a sum over repeated indices is understood. In the limit \( \lambda_a u_a \ll M \), the singlets \( \nu_j \) acquire a mass given by

\[
m = (\lambda_{\alpha}^T u_a) M^{-1} (\lambda_a u_a) \ll M,
\]

while the mass matrix of the three light neutrinos is

\[
m_L = h m_{\nu}^{-1} h^T v^2.
\]

Therefore the smallness of the \( \nu_j \) masses is due to a seesaw mechanism involving just the SM singlet leptons.

There are different variants of this scenario, depending on whether lepton number is a global or local symmetry, and the time of spontaneous breaking, denoted by the temperature \( T_L \). We first consider the case of global \( U(1)_L \). If lepton number is broken after the DM freeze out, WB does not work because the \( \chi \) field, being charged, can also hold an asymmetry. In turn this asymmetry induces a washout proportional to \( \frac{\gamma_{(\chi \chi \rightarrow N \nu)}}{n_{\chi} H(z)} \), which freezes out at the same moment as the annihilation of DM, violating one of the basic requirements of WB [14].

However, it may be possible that WB occurs via helicogenesis when \( U(1)_L \) is already broken, i.e., \( T_L > T_h \) or \( T_L > T_{sfo} \), being \( T_h \) the temperature at which the \( \nu_j \) helicity asymmetry is generated. If \( T_L \sim \) few TeV, it is natural that the physical masses of the singlet scalars \( S_a \) after \( U(1)_L \)

6 The baryon asymmetry is roughly given \( Y_B \sim \frac{1}{2} [Y_\chi(z_w) - Y_\chi(\infty)] \), where \( z_w \) is the value of \( z \) at which washouts freeze out and \( Y_\chi(\infty) \) is the relic DM density normalized to the entropy density. Given that \( \Omega_{DM} \sim 5 \Omega_B \) and \( m_\chi \sim 1 \) TeV, it is clear that all washout processes must freeze out before DM annihilations, when \( Y_\chi \) is several orders of magnitude above its final value \( Y_\chi(\infty) \).
breaking satisfy $m_{S_{A}} > m_{\chi_{1}} \gtrsim 1 \text{ TeV}$ in a broad region of the parameter space. The sterile neutrinos $\nu_{j}$ acquire a mass given by Eq. (18), therefore the condition $m_{j} \ll T_{sfo} \sim 100 \text{ GeV}$ needed to generate the helicity asymmetry in $\nu_{j}$, leads to $(\lambda_{a} u_{a})^{2}/M \ll 100 \text{ GeV}$. For instance, assuming $u_{a} \sim M = 1 \text{ TeV}$, $\lambda_{a}$ of order 0.2 is required to obtain $m_{j} \sim 40 \text{ GeV}$. On the other hand, the Yukawa couplings $\lambda_{a \nu_{i} j}$ should be sizeable, of $\mathcal{O}(1)$, to have enough CP violation and get the correct DM relic abundance, so there is some tension between these two requirements for WB. Since the neutrino masses depend only on the couplings $\lambda_{a}$ while a combination of $\lambda_{a}$ and $\nu_{a}$ appears in the CP-asymmetry, it is conceivable that some cancellations due to phases allow to satisfy all constraints in certain regions of the parameter space with only two singlet scalars. Alternatively, in the presence of three scalars it may happen that $u_{1}, u_{2} \ll u_{3}$ but $\lambda_{1}, \lambda_{2} \gg \lambda_{3 j j}$, achieving a large enough CP asymmetry through the couplings $\lambda_{1,2}$ and getting $m_{j} \ll T_{sfo}$ without accidental cancellations. Moreover, the vev’s at $T_{h} \gg T_{sfo}$ may be different from the vev’s at $T = 0$, as in the singlet Majoron model [36], helping to enlarge the allowed parameter space. We thus conclude that WB seems feasible in the present framework if $T_{L} > T_{h} > T_{sfo}$.

The spontaneous breaking of a global symmetry leads to a massless Goldstone boson, the Majoron,

$$J = \sum_{a} \frac{u_{a}}{u} \text{Im} S_{a}, \quad (20)$$

with $u = \sqrt{\sum_{a} u_{a}^{2}}$. However, non-perturbative gravitational effects are expected to explicitly break global symmetries and provide a mass to the Majoron [37]. If this mass is $m_{J} \lesssim \text{ few hundred GeV}$, processes mediated by $N$ such as $\nu_{1}^{+} J \rightarrow \nu_{2}^{-} J$ could lead to a fast washout of the $\nu$ helicity asymmetry and a more detailed analysis is required.

This potential problem can be avoided by promoting lepton number to a gauge symmetry. In this framework, only the case $U(1)_{B-L}$ is anomaly free without requiring new exotic fermions to cancel anomalies. Then, there are additional constraints due to the searches of the extra $Z'$ gauge boson at LEP, Tevatron and LHC. While LHC searches for heavy resonances depend both on the $U(1)_{B-L}$ coupling strength $g_{B-L}$ and $Z'$ mass [38, 39], limits from LEP II imply a model independent bound on the vev $u = M_{Z'}/(2 g_{B-L}) \gtrsim 3 \text{ TeV}$ [40]. Thus in the gauged case $U(1)_{B-L}$ is necessarily broken before the electroweak phase transition, and the results discussed above for $T_{L} > T_{h} > T_{sfo}$ apply. Now there is no Majoron, and the only concern would be that the cross sections of the new lepton-number-conserving annihilation channels induced by the $Z'$ boson are not too large, so a significant fraction of DM annihilations still proceed through lepton-number-violating processes with $\nu$ in the final state, leading to a sizeable helicity asymmetry in the $\nu$ population. This requirement is easy to satisfy, since $g_{B-L}$ as well as $M_{Z'}$ are free parameters.

Notice that the suppression of fast washouts is also at work when the helicity asymmetry in the “light” sterile neutrinos $\nu$ is generated during the out-of-equilibrium decay of the heavy states, namely the singlet fermions $N$ or the singlet scalars. In this case, the decaying particle should have tiny Yukawa couplings, to achieve the out-of-equilibrium condition at low temperatures, $T \sim \mathcal{O} \text{ TeV}$, and the direct connection between the baryon asymmetry and the DM relic abundance is lost [26].

4. Summary and outlook

Thermal baryogenesis at low scales, $\mathcal{O}$(TeV), suffers from the intrinsic problem of fast asymmetry washout. We have proposed a mechanism which overcomes this difficulty: massive decay or annihilation products which do not store asymmetry, such as Majorana fermions or real scalars. We have illustrated the mechanism in the case of leptogenesis via DM annihilation into sterile neutrinos: one of these neutrinos, $N$, must be ”heavy” ($M \sim \mathcal{O}$ TeV), providing a Boltzmann distribution with $\mathcal{O}(1)$.

Note that if the $Z'$ can decay into sterile neutrinos, the LHC limits may be relaxed [41].
suppression of the washouts. The other neutrino, $\nu$, should be relatively "light" ($m_\nu \ll 100$ GeV), so its mass is negligible during the annihilation epoch, allowing for an helicity asymmetry to be generated (helicitogenesis). This asymmetry is partially transferred to SM leptons via Yukawa interactions and subsequently reprocessed into a baryon asymmetry via the electroweak spallation.

At least one of the sterile neutrino should have $m_\nu \sim \mu$ GeV and some of the Yukawa couplings must be larger than $2 \times 10^{-7}$, therefore the mass of at least one of the SM neutrinos results larger than few $\times 0.01$ eV, barring accidental cancellations, in the ballpark of the mass scales set by neutrino oscillations. Therefore, helicitogenesis from the annihilation of DM connects three open problems of particle physics, namely the BAU, DM and neutrino masses.

The required pattern of sterile neutrino masses and couplings appears naturally in the so-called double seesaw scenario, where the smallness of the $\nu$ masses can be due to a $U(1)_L$ symmetry spontaneously broken. Thus, we have constructed an extended $U(1)_L$ symmetric double-seesaw model, including also fermionic DM and two SM singlet scalars, all of them charged under lepton number. We have shown that it is possible to reconcile the helicitogenesis requirements with the measured light neutrino masses provided that $U(1)_L$ breaks spontaneously prior to DM freeze out, or heavy particle decay, and hence before the electroweak phase transition. Within this framework, it seems feasible to have successful WB and explain the observed light neutrino masses with DM and sterile neutrino couplings to the singlet scalars close to $O(1)$. The presence of a light Majoron associated to the breaking of the global $U(1)_L$ symmetry is a potential problem, which can be solved by gauging $U(1)_{B-L}$.

From the phenomenological point of view, the models that we have discussed involve new particles at the TeV scale, so in principle they can be tested in current or near-future experiments. Especially interesting would be to analyze whether the prospects for detecting the sterile neutrinos, generically very difficult in Type I seesaw models, are improved by their additional interactions with the singlet scalars, which in general mix with the SM Higgs field, or with the $Z'$ boson in the $U(1)_{B-L}$ gauged case. With respect to DM detection, the phenomenology of WB models has been extensively analyzed neglecting the mixing among the SM Higgs and the extra scalar singlets [14, 24]. However the impact of the mixing within the scalar sector and of the $Z'$ interaction deserves further investigation.

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