

Low-scale leptogenesis in the scotogenic model: Spectator processes and benchmark points

J. Racker^{*}

Instituto de Astronomía Teórica y Experimental (IATE), Consejo Nacional de Investigaciones Científicas y Técnicas (CONICET)—Universidad Nacional de Córdoba (UNC), Laprida 854, X5000BGR, Córdoba, Argentina and Observatorio Astronómico de Córdoba (OAC), Universidad Nacional de Córdoba (UNC), Laprida 854, X5000BGR, Córdoba, Argentina



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We study leptogenesis from the decay of the lightest sterile neutrino in the scotogenic model with a scalar dark matter candidate. Our analysis focuses on the possible exponential suppression of washouts for sizable values of the inert Higgs mass and the crucial role of some spectator processes for this to happen. We show that leptogenesis can be successful for TeV-scale masses of the lightest sterile neutrino for both, thermal and zero, initial abundances of the neutrinos, without resorting to quasidegenerate neutrinos. Moreover, leptogenesis is viable for Yukawa couplings of the heavier sterile neutrinos large enough to yield observable charged lepton flavor violation (CLFV) processes in current and planned experiments.

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Introduction. The baryon density of the Universe has been determined to better than one percent by *Planck* [1] and is consistent with the Big-Bang nucleosynthesis prediction of light element abundances (see, e.g., [2]). Given that the Universe does not seem to contain significant amounts of antimatter [3], the baryon density is also a measure of the baryon asymmetry of the Universe (BAU), which, normalized to the entropy density, s , is equal to $Y_B \equiv n_B/s \simeq 8.7 \times 10^{-11}$ [1]. The BAU is known to be one of the strongest evidences of physics beyond the Standard Model (SM) [4–7].

Baryogenesis via leptogenesis in neutrino mass models [8], notably in the type I seesaw, is one of the best motivated explanations for the BAU, but unfortunately these scenarios typically require masses of the exotic particles too high to allow for experimental exploration in foreseeable experiments (beyond the connection with light neutrino masses). In the type I seesaw with hierarchical heavy neutrino masses, the Davidson-Ibarra bound [9] for the CP asymmetry in neutrino decays implies that the mass of the lightest heavy neutrino (M) must satisfy $M \gtrsim 10^8\text{--}10^9$ GeV. More generally, low-scale leptogenesis generally suffers from strong washout processes that erase the asymmetry. Still, it has also been realized that low-scale

leptogenesis can be successful in particular regions of the parameter space of the type I seesaw model [10–13], including regions with measurable active-sterile mixings (see [14] for a review and references, and, e.g., [15–18] for the connection between leptogenesis via neutrino oscillations and resonant leptogenesis).

Remarkably, the type I seesaw with one KeV-scale neutrino may also provide a dark matter (DM) candidate, however the correct amount of BAU and DM can only be obtained in outstandingly fine tuned regions of the parameter space [11,19–21]. The phenomenology changes substantially if, together with the sterile neutrinos, a second Higgs doublet is added to the SM particle content, and the new particles are required to be odd under a discrete Z_2 symmetry (while the SM fields are even). This is the scotogenic model [22,23], which is arguably the simplest and most studied model to explain and relate the smallness of neutrino masses and DM.

Soon after the publication of [23] it was realized that baryogenesis via leptogenesis is also possible in this model and that the lower bound on M due to the Davidson-Ibarra bound could be relaxed [24,25]. However, only estimates were presented for the crucial washout effects. Indeed more detailed treatments including all neutrino oscillation data were performed in [26,27], concluding that successful leptogenesis with TeV-scale masses seemed possible only for quasidegenerate neutrinos resorting to the resonant leptogenesis mechanism. A more exhaustive exploration of the parameter space for leptogenesis in the scotogenic model was performed in [28,29]. The results of [28] suggested that nonresonant leptogenesis at the TeV-scale from the decay of the lightest sterile neutrino was feasible,

^{*}Contact author: jracker@unc.edu.ar

either by assuming an initial thermal abundance for this species and/or for sizable masses of the inert Higgs (the new Higgs doublet) in its role of being one of the decay products of the sterile neutrinos. However, the work [28] centered on the mechanisms to avoid the strong washouts and it did not attempt to fit neutrino oscillation data. Moreover, it performed a simplified treatment of some crucial spectator processes (i.e., processes which, while not being the source of the lepton asymmetry, they are active during leptogenesis and influence the resultant BAU). Instead, in [29] the exploration of the parameter space was performed accounting for all neutrino oscillation data (via a Casas-Ibarra [30] parametrization of the Yukawa couplings), but the flavor structure in the transport equations [31–38] was neglected and, more importantly, the effects of the inert Higgs mass on leptogenesis were not considered. The lower bound on the mass of the lightest sterile neutrino, assuming a thermal initial abundance, was found to be around 10 TeV [39].

Here we use an extended set of flavored Boltzmann equations (BEs), including explicitly some crucial spectator processes, to show that for sizable values of the inert Higgs mass, the parameter space compatible with DM, neutrino masses and successful leptogenesis opens up considerably, embracing regions which are experimentally interesting.

The model and Boltzmann equations. The scotogenic model extends the SM by a second Higgs doublet η and (two or) three singlet fermions N_i (aka sterile neutrinos) with masses M_i . The model also has a discrete Z_2 symmetry under which the new fields are odd and the SM fields are even. The scalar potential and Yukawa terms in the Lagrangian are given by

$$-V(\phi, \eta) - (h_{ai} \bar{N}_i \tilde{\eta}^\dagger \ell_\alpha + \text{H.c.}) \subset \mathcal{L},$$

where $\alpha = e, \mu, \tau$ and $i = 1, 2, 3$ are flavor indices, ℓ_α are the leptonic $SU(2)$ doublets, $\eta = (\eta^+, (\eta_R + i\eta_I)/\sqrt{2})^T$ is the inert Higgs ($\tilde{\eta} = i\sigma_2 \eta^*$), ϕ is the SM Higgs and the scalar potential $V(\phi, \eta)$ can be written in terms of real parameters $m_\phi^2, m_\eta^2, \lambda_{1,2,3,4,5}$ as

$$\begin{aligned} V(\phi, \eta) = & m_\phi^2 \phi^\dagger \phi + m_\eta^2 \eta^\dagger \eta + \frac{\lambda_1}{2} (\phi^\dagger \phi)^2 + \frac{\lambda_2}{2} (\eta^\dagger \eta)^2 \\ & + \lambda_3 (\phi^\dagger \phi)(\eta^\dagger \eta) + \lambda_4 (\eta^\dagger \phi)(\phi^\dagger \eta) \\ & + \frac{\lambda_5}{2} [(\eta^\dagger \phi)^2 + (\phi^\dagger \eta)^2]. \end{aligned}$$

After electroweak symmetry breaking, the new physical scalar states have masses

$$\begin{aligned} m_{\eta_R}^2 &= m_\eta^2 + (\lambda_3 + \lambda_4 + \lambda_5) v^2, & m_{\eta^\pm}^2 &= m_\eta^2 + \lambda_3 v^2, \\ m_{\eta_I}^2 &= m_\eta^2 + (\lambda_3 + \lambda_4 - \lambda_5) v^2, \end{aligned} \quad (1)$$

with $v = 246/\sqrt{2}$ GeV the vacuum expectation value of the Higgs.

Here we will consider the case that one of the neutral components of η is the lightest -stable- odd particle under the Z_2 symmetry. For definiteness we take $\lambda_5 > 0$ so that η_I is the DM candidate and restrict our study to the mass window $530 \text{ GeV} \lesssim m_{\eta_I} \lesssim 20 \text{ TeV}$, where the correct relic abundance can be obtained with a proper choice of the scalar couplings (for bounds and related studies see, e.g., [25,46–62]).

Neutrino masses are generated radiatively (due to the Z_2 symmetry) and they are also suppressed by λ_5 , which is a natural parameter in the sense of 't Hooft, as a conserved lepton number can be defined in the limit $\lambda_5 \rightarrow 0$. At one-loop order the mass matrix of the active neutrinos is given by [23,63]

$$(m_\nu)_{\alpha\beta} = (h\Lambda h^T)_{\alpha\beta}, \quad (2)$$

where Λ is a diagonal matrix with entries

$$\Lambda_{ii} = \frac{M_i}{32\pi^2} \left[\frac{m_{\eta_R}^2}{m_{\eta_R}^2 - M_i^2} \ln \frac{m_{\eta_R}^2}{M_i^2} - \frac{m_{\eta_I}^2}{m_{\eta_I}^2 - M_i^2} \ln \frac{m_{\eta_I}^2}{M_i^2} \right]. \quad (3)$$

The mass matrix m_ν can be diagonalized by the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix, U , to yield the light neutrino masses, $U^T m_\nu U = d_\nu = \text{diag}(m_1, m_2, m_3)$. From these relations it is possible to obtain a Casas-Ibarra [30] parametrization for the Yukawa coupling matrix $h = (h_{ai})$ in terms of a complex orthogonal matrix R ,

$$h = U^* \sqrt{d_\nu} R^T \sqrt{\Lambda}^{-1}. \quad (4)$$

In turn, the matrix R can be written as the product of three rotations by complex angles z_{23} , z_{13} and z_{12} , $R = R_{23} R_{13} R_{12}$. Below, each complex angle will be decomposed into its real and imaginary parts, namely $z_{ij} = z_{ij_R} + z_{ij_I} i$. For our numerical study we will calculate the Yukawa couplings from this parametrization, taking for U and the neutrino mass-squared differences the expression and central values given in [64]. The only free parameters remaining in U will be the two Majorana phases, $\eta_{1,2}$.

As other models for neutrino masses, the scotogenic model also has the virtue of incorporating a baryogenesis mechanism [8]. Namely, as the sterile neutrinos decay in the primitive Universe, generally some amount of lepton asymmetry is produced, which is partially converted into a baryon asymmetry by the electroweak sphaleron processes [65]. The baryon asymmetry freezes at a temperature $T \sim 130 \text{ GeV}$ [66], when the Hubble rate becomes faster than the sphaleron processes.

Baryogenesis can be described by a set of BEs. At the low temperatures we are interested, all charged lepton Yukawa interactions are fast and the evolution of the lepton

asymmetry can be followed appropriately by one BE for each lepton flavor (see the references on flavor quoted above). It also becomes necessary to trace the sterile neutrino abundances and, finally, we will add a BE to track the evolution of the asymmetry in the inert Higgs field, which plays a key role as we will show. Under several

-reasonable and common- approximations, including kinetic (but not chemical) equilibrium of all particle species, Maxwell-Boltzmann statistics, neglect of several scattering processes and working at linear order in the asymmetries (see, e.g., [67–71]), an appropriate set of BEs for our study is

$$\begin{aligned}
 \frac{dY_{N_i}}{dz} &= \frac{-1}{sH\dot{z}} \left(\frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} - 1 \right) \sum_{\alpha} 2\gamma_{\ell_{\alpha}\eta}^{N_i}, \\
 \frac{dY_{\Delta_{\alpha}}}{dz} &= \frac{-1}{sH\dot{z}} \left\{ \sum_i \left(\frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} - 1 \right) \epsilon_{ai} 2\gamma_{\ell_{\alpha}\eta}^{N_i} - \sum_i \gamma_{\ell_{\alpha}\eta}^{N_i} [y_{\ell_{\alpha}} + y_{\eta}] \right. \\
 &\quad \left. - \sum_{\beta} (\gamma_{\ell_{\beta}\eta}^{\ell_{\alpha}\eta'} + (1 + \delta_{\alpha\beta}) \gamma_{\eta\bar{\eta}}^{\ell_{\alpha}\ell_{\beta}}) [y_{\ell_{\alpha}} + y_{\ell_{\beta}} + 2y_{\eta}] - \sum_{\beta \neq \alpha} (\gamma_{\ell_{\beta}\eta}^{\ell_{\alpha}\eta'} + \gamma_{\ell_{\beta}\eta}^{\ell_{\alpha}\bar{\eta}} + \gamma_{\eta\bar{\eta}}^{\ell_{\alpha}\ell_{\beta}}) [y_{\ell_{\alpha}} - y_{\ell_{\beta}}] \right\}, \\
 \frac{dY_{\Delta_{\eta}}}{dz} &= \frac{1}{sH\dot{z}} \left\{ \sum_{i,\alpha} \left(\frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} - 1 \right) \epsilon_{ai} 2\gamma_{\ell_{\alpha}\eta}^{N_i} - \sum_{i,\alpha} \gamma_{\ell_{\alpha}\eta}^{N_i} [y_{\ell_{\alpha}} + y_{\eta}] - \sum_{\alpha,\beta} (\gamma_{\ell_{\beta}\eta}^{\ell_{\alpha}\eta'} + 2\gamma_{\eta\bar{\eta}}^{\ell_{\alpha}\ell_{\beta}}) [y_{\ell_{\alpha}} + y_{\ell_{\beta}} + 2y_{\eta}] \right. \\
 &\quad \left. - 2(\gamma_{\phi\eta}^{\bar{\phi}\eta} + 2\gamma_{\phi\phi}^{\eta\eta}) [y_{\eta} - y_{\phi}] \right\}, \tag{5}
 \end{aligned}$$

with $i = 1, 2, 3$, $\alpha = e, \mu, \tau$, $z \equiv M_1/T$, H is the Hubble rate, $Y_X \equiv n_X/s$ is the number density of particle X normalized to the entropy density, and $y_X \equiv Y_{\Delta_X}/Y_X^{\text{eq}} \equiv (Y_X - Y_{\bar{X}})/Y_X^{\text{eq}}$ is the corresponding number density asymmetry normalized to the equilibrium density. We have also defined $Y_{\Delta_{\alpha}} \equiv Y_B/3 - Y_{L_{\alpha}}$, where Y_B is the baryon asymmetry defined previously and $Y_{L_{\alpha}}$ is the total lepton asymmetry in flavor α normalized to the entropy density. We use the notation γ_f^i for the reaction density (i.e., rate per unit volume) of the corresponding process $i \rightarrow f$. The prime on some rates indicates that the contribution from on-shell neutrinos has to be subtracted [67,69], and the term with the Kronecker delta ($\delta_{\alpha\beta}$) takes into account that processes with $\alpha = \beta$ change $Y_{\Delta_{\alpha}}$ by two units. Finally, the flavored CP asymmetries in N_i decays are defined by $\epsilon_{ai} = (\gamma_{\ell_{\alpha}\eta}^{N_i} - \gamma_{\ell_{\alpha}\bar{\eta}}^{N_i}) / \sum_{\beta} (\gamma_{\ell_{\beta}\eta}^{N_i} + \gamma_{\ell_{\beta}\bar{\eta}}^{N_i})$ [72].

The set of BEs can be closed by additional relations among the number density asymmetries, which are obtained considering fast interactions (that lead to equalities among chemical potentials) and the conserved charges [74]. Performing an analysis quite similar to the lowest temperature regime described in [36], but including the additional scalar doublet η , we obtain

$$\begin{aligned}
 Y_{\Delta_{\ell_{\alpha}}} &= [32(Y_{B-L} - Y_{\Delta_{\alpha}}) - 442Y_{\Delta_{\alpha}} - 27Y_{\Delta_{\eta}}]/1422, \\
 Y_{\Delta_{\phi}} &= -(8Y_{B-L} + 13Y_{\Delta_{\eta}})/79, \quad \text{and} \\
 Y_B &= (28Y_{B-L} + 6Y_{\Delta_{\eta}})/79, \tag{6}
 \end{aligned}$$

with $Y_{B-L} \equiv \sum_{\alpha} Y_{\Delta_{\alpha}}$.

Spectator processes and benchmark points. We have made a rough exploration of the parameter space compatible with the data on neutrino mixing, successful leptogenesis and a scalar DM candidate with $m_{\eta} \gtrsim 530$ GeV [75]. Notice that the rates relevant for leptogenesis depend on the mass of the inert Higgs before the electroweak symmetry breaking (m_{η}), while the masses of the new scalar states after the electroweak symmetry breaking are given by Eq. (1), where the quartic couplings can be chosen to obtain the observed DM relic density (see, for instance, the Appendix C of [56]). Below we describe four benchmark points (BPs) that uncover interesting regions of parameter space and discuss the role of some spectator processes. The values of the parameters are given in Table I and the moduli of the Yukawa couplings in Table II. Note that the small values of λ_5 are motivated by the requirement of having sizable Yukawa couplings, but they are large enough to satisfy DM direct detection bounds [25,46].

BP1: This BP shows that leptogenesis in the scotogenic model is possible for masses as low as $M_1 \sim 1.3$ TeV, assuming a thermal initial abundance of N_1 , i.e., almost one order of magnitude below the lower bound found in [29]. We stress that lowering the mass scale of (nonresonant) leptogenesis one order of magnitude, particularly at these already low energy scales for thermal leptogenesis, is not trivial. In the scotogenic model this can be achieved thanks to the sizable mass of one of the states coupled to the sterile neutrinos, namely the inert Higgs η , so that the washouts mediated by the sterile neutrinos are Boltzmann suppressed for $T \lesssim m_{\eta}$. This way of suppressing washouts can be very effective, but for a proper treatment it is crucial to take into account that in general the inert Higgs is not in chemical

TABLE I. Benchmark points for leptogenesis in the scotogenic model. The Yukawa couplings have been obtained from the Casas-Ibarra parametrization given in Eq. (4), with $z_{12_R} = z_{12_l} = z_{23_R} = 0$ for all points. Neutrino mixing data has been taken equal to the central values reported in [64] assuming a normal hierarchy. To evaluate Λ we have taken $m_{\eta_l}^2 \simeq m_{\eta}^2$ and $m_{\eta_R}^2 = m_{\eta_l}^2 + 2\lambda_5 v^2$ in Eq. (3). Sphalerons freeze out at $T = 130$ GeV and for all points $Y_B \simeq 8.7 \times 10^{-11}$.

BP	M_1 [TeV]	M_2 [TeV]	M_3 [TeV]	m_η [TeV]	λ_5	z_{13_R}	z_{13_l}	z_{23_l}	η_1	η_2	m_1 [eV]
1	1.3	6	12	0.6	1.4×10^{-5}	9.5×10^{-6}	9.5×10^{-6}	0	$\pi/2$	0	10^{-15}
2 ^a	3	6	9	0.9	1×10^{-5}	3×10^{-6}	-3×10^{-6}	2.5	1.2	0	10^{-13}
3 ^b	4	5.5	8	1.6	1.23×10^{-4}	6.5×10^{-6}	-6.5×10^{-6}	6	$\pi/2$	0	10^{-12}
4 ^c	2.5	$\sqrt{10}M_1$	$10M_1$	$0.7M_1$	1×10^{-3}	6×10^{-5}	-6×10^{-5}	5	$\pi/2$	$\pi/4$	10^{-11}

^a $\text{CR}(\mu - e, \text{Ti}) = 1.3 \times 10^{-19}$.

^b $\text{Br}(\mu \rightarrow e\gamma) = 5 \times 10^{-13}$. $\text{CR}(\mu - e, \text{Ti}) = 2.1 \times 10^{-15}$, $\text{Br}(\mu \rightarrow 3e) = 3 \times 10^{-15}$ (dipole contributions).

^cZero initial abundance of N_1 , $\text{CR}(\mu - e, \text{Ti}) = 9.3 \times 10^{-20}$.

TABLE II. Moduli of the Yukawa couplings for the benchmark points of Table I, rounded to 2 significant figures.

BP	$ h_{e1} $	$ h_{\mu 1} $	$ h_{\tau 1} $	$ h_{e2} $	$ h_{\mu 2} $	$ h_{\tau 2} $	$ h_{e3} $	$ h_{\mu 3} $	$ h_{\tau 3} $
1	9.4×10^{-9}	4.4×10^{-8}	3.9×10^{-8}	1.3×10^{-3}	1.4×10^{-3}	1.3×10^{-3}	9.9×10^{-4}	4.9×10^{-3}	4.4×10^{-3}
2	6.5×10^{-9}	1.8×10^{-8}	2.2×10^{-8}	1.5×10^{-2}	3.4×10^{-2}	3.1×10^{-2}	1.6×10^{-2}	3.8×10^{-2}	3.4×10^{-2}
3	9×10^{-9}	1.4×10^{-8}	1.9×10^{-8}	1.7×10^{-1}	3.9×10^{-1}	3.5×10^{-1}	1.8×10^{-1}	4.1×10^{-1}	3.7×10^{-1}
4	1.4×10^{-8}	5.4×10^{-8}	5.5×10^{-8}	2.6×10^{-2}	4.1×10^{-2}	5.9×10^{-2}	3.4×10^{-2}	5.2×10^{-2}	7.6×10^{-2}

equilibrium and has a nonzero chemical potential (note in particular that a number density asymmetry of η particles is also generated during decays). The η asymmetry influences the evolution of the lepton asymmetries. This is because an excess number of, e.g., η over $\bar{\eta}$, also induces washout of the lepton asymmetry, as it increases the rate of, e.g., $\ell_\alpha \eta \rightarrow \bar{\ell}_\beta \bar{\eta}$ over $\bar{\ell}_\alpha \bar{\eta} \rightarrow \ell_\beta \eta$. The evolution of this asymmetry ($Y_{\Delta\eta}$) is regulated by the rates of the spectator processes $\bar{\phi}\eta \rightarrow \phi\bar{\eta}$ and $\eta\eta \rightarrow \phi\phi$, which are proportional to λ_5^2 [see the last equation of the set (5)]. If fast enough, they lead to an exponential depletion of the η -asymmetry, but if not, $Y_{\Delta\eta}$ is proportional to the lepton asymmetries. In the latter case, washout terms in the BEs like $\gamma_{\bar{\ell}_\beta \bar{\eta}}^{\ell_\alpha \eta} y_\eta$ are not exponentially suppressed (recall that $y_\eta = Y_{\Delta\eta}/Y_\eta^{\text{eq}}$).

To show the importance of these spectator processes, we have solved the BEs (5) for BP1, but multiplying the rates of $\bar{\phi}\eta \rightarrow \phi\bar{\eta}$ and $\eta\eta \rightarrow \phi\phi$ by an arbitrary parameter λ . The resultant baryon asymmetry as a function of λ is represented by the blue curve in Fig. 1. In the limit $\lambda \rightarrow \infty$ spectator processes are faster than the Hubble rate during the whole leptogenesis epoch, ensuring the equality among the chemical potentials of the inert and SM Higgs fields, $\mu_\eta = \mu_\phi$, which implies a Boltzmann suppression of $Y_{\Delta\eta}$ relative to $Y_{\Delta\phi}$ for $T \lesssim m_\eta$. Instead, for $\lambda \rightarrow 0$ these spectator processes are absent, which effectively corresponds to the conservation of a modified lepton number, implying that $Y_{\Delta\eta}$ must be equal to a linear combination of $Y_{\Delta_{e,\mu,\tau}}$ and the corresponding washouts are not exponentially suppressed. For this BP there is a difference of more than one order of magnitude between the two limits and the

actual asymmetry, i.e., the one with $\lambda = 1$, is close to the maximum value obtained when $\mu_\eta = \mu_\phi$. We note that this effect was discussed only qualitatively in [28] (which assumed the equality $\mu_\eta = \mu_\phi$). It was also presented more quantitatively in a proceedings [76], but without any details on the BEs and in a toy scenario (not fitting neutrino data).

In some regions of parameter space, the simplifying assumption $\mu_\eta = \mu_\phi$ (that allows to reduce the set of BEs) can lead to a considerable overestimate of the baryon asymmetry. This is illustrated in Fig. 2, where the baryon asymmetry obtained with this assumption is plotted as a function of z_{13_R} and compared to the more precise calculation using the full set of BEs (5). For these points of parameter space Y_B is overestimated by a factor of 2 to 3, giving in particular -wrong- values above the observed BAU for $3 \times 10^{-6} < z_{13_R} < 9 \times 10^{-6}$ [77].

BP2: The first BP is interesting because it shows that—and how—leptogenesis is possible with very low values of the sterile neutrino masses in one of the simplest extensions of the SM with a natural DM candidate. However, the Yukawa couplings of all neutrino species are too tiny for them to leave signals in current or planned experiments. BP2 reveals that successful leptogenesis is possible with higher values of the Yukawa couplings of N_2 and N_3 , large enough to induce CLFV processes that could be observed in the foreseeable future. More specifically, for this BP we obtain that the muon-to-electron conversion rate in titanium is larger than 10^{-19} (relative to the muon capture rate), above the sensitivity that could be achieved in planned experiments [80]. We have calculated the conversion rate via the approximation

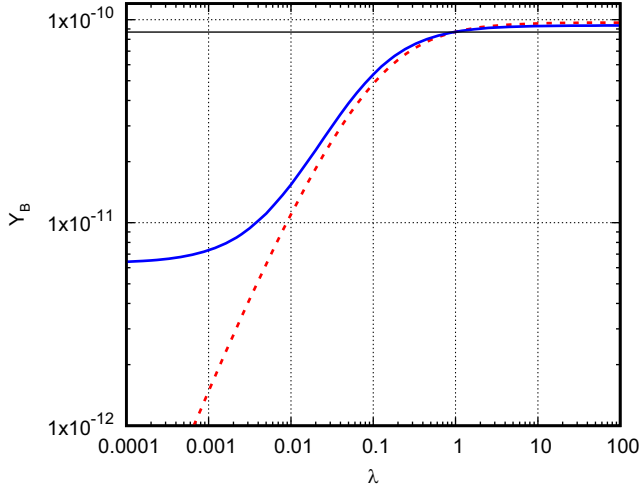


FIG. 1. The baryon asymmetry obtained for BP1 (solid blue curve) and BP4 (dashed red curve) integrating the BEs (5), but with the spectator processes reaction densities $\gamma_{\phi\eta}^{\phi\eta}$ and $\gamma_{\phi\phi}^{\eta\eta}$ multiplied by λ , as a function of this parameter λ . For reference, the observed value of Y_B is depicted with the horizontal solid black line.

$$\text{CR}(\mu - e, \text{Ti}) \approx \text{Br}(\mu \rightarrow e\gamma) B(A, Z)/428, \quad (7)$$

with $B(A, Z) = 1.8$ [81] (the lowest value quoted in that work), valid for a dominant dipole contribution. In turn, the branching ratio $\text{Br}(\mu \rightarrow e\gamma)$ has been computed using the analytical expression of [82], taking $m_{\eta^\pm}^2 \simeq m_\eta^2$.

For this BP the imaginary part of the complex angle z_{23} is moderately large ($z_{23_I} = 2.5$), which results in Yukawa couplings larger by a factor $\sim \cosh z_{23_I}$ than a rough estimate from Eq. (2) disregarding complex phases. Therefore, for experimental probes the regions of parameter space with large imaginary parts of complex angles are the most interesting ones. They may be motivated by symmetries, but also renormalization group running effects might be important (see [63] and [83] for a different parametrization of the orthogonal matrix R motivated by symmetry considerations). In any case, here we have shown that leptogenesis in the scotogenic model is possible for Yukawa couplings large enough to allow for experimental exploration and in regions of parameter space different from those allowed by resonant leptogenesis or leptogenesis via neutrino oscillations.

BP3: The Boltzmann suppression of washouts due to the mass of the inert Higgs can be very effective, enabling successful leptogenesis for Yukawa couplings much larger than those of BP2. For BP3 $z_{23_I} = 6$ and the $\text{Br}(\mu \rightarrow e\gamma) = 5 \times 10^{-13}$, which is above the current bound, $\text{Br}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}$ (90% confidence level) [84]. Therefore this BP shows that some regions of parameter space compatible with successful leptogenesis are already ruled out by experiments. The $\text{CR}(\mu - e, \text{Ti})$ obtained from Eq. (7) is below the current bound, but much larger than

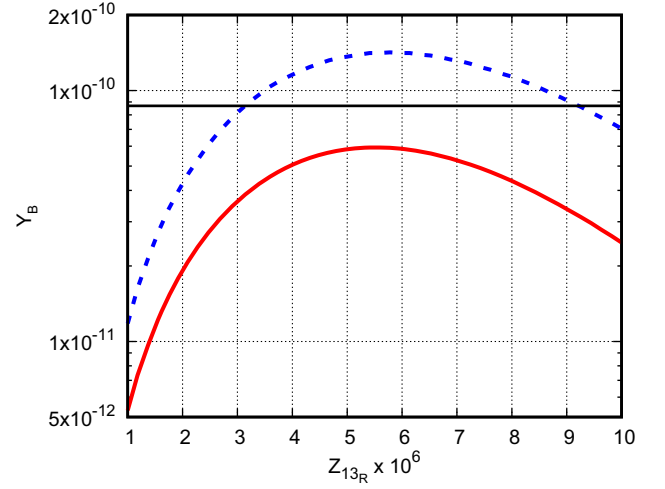


FIG. 2. The baryon asymmetry as a function of z_{13_R} assuming $\mu_\phi = \mu_\eta$ (blue-dashed curve) and using the complete set of BEs (5) (red-solid curve). Here we have taken $M_1 = 1.6$ TeV, $M_2 = \sqrt{10}M_1$, $M_3 = 10M_1$, $m_\eta = 0.6M_1$, $\lambda_5 = 8 \times 10^{-6}$, $\eta_1 = \pi/2$, $\eta_2 = 0$, $m_1 = 10^{-13}$ eV, $z_{13_I} = z_{13_R}$ and $z_{12} = z_{23} = 0$. Note that $z_{13_R} \in [10^{-6}, 10^{-5}]$ and that we have chosen z_{13_R} as the independent parameter (with $z_{13_I} = z_{13_R}$) just to illustrate a continuous set of points in the parameter space where the baryon asymmetry varies substantially around the observed value (which is mainly due to the change in the decay width of N_1). The observed value of Y_B is depicted with the horizontal solid black line.

sensitivities expected in planned experiments. Actually, this expression for the $\text{CR}(\mu - e, \text{Ti})$ only takes into account the dipole contribution, but there is also a nondipole contribution that is dominant in some regions of the parameter space of the scotogenic model [82], leading to much larger values of the $\mu - e$ conversion rate than the obtained from Eq. (7). Given that successful leptogenesis is possible for such large Yukawa couplings, the possibility opens to explore regions of parameter space where the ratio $\text{CR}(\mu - e, \text{Ti})/\text{Br}(\mu \rightarrow e\gamma)$ departs from the relation given in Eq. (7) and becomes a more distinct signal of the scotogenic model. Similar considerations apply to the ratio $\text{Br}(\mu \rightarrow 3e)/\text{Br}(\mu \rightarrow e\gamma)$, given that the $\text{Br}(\mu \rightarrow 3e)$ also is above the sensitivity of planned experiments for this BP (see [82]).

BP4: For the previous BPs we have assumed a thermal initial abundance of the singlet fermion N_1 (as was also assumed in [29] and in many other models of low-scale baryogenesis from particle decays). This amounts to assume that the sterile neutrinos were produced at higher temperatures by some other—unknown—process not related to the Yukawa interactions of the scotogenic model. With BP4 we show that, perhaps unexpectedly, low-scale leptogenesis with $M_1 \sim 2.5$ TeV is also possible starting with a zero initial abundance of N_1 , i.e., with neutrino production being due just to the Yukawa interactions of the model. This requires large values of z_{23_I} , and for this BP the

$\text{CR}(\mu - e, \text{Ti})$ lies within the limit of the sensitivity of planned experiments. The spectator processes are even more important here, as can be seen from the red curve in Fig. 1, which has been obtained in an analogous way to the blue curve for BP1. Note that, although λ_5 is two orders of magnitude larger than for BP1, the baryon asymmetry decreases by one order of magnitude for $\lambda \sim 0.01$ and by several orders of magnitude if the spectator processes are turned off.

As a final comment, we note that for all the BPs above, the electric dipole moment of the electron estimated from [85] is -at least- a few orders of magnitude below expected sensitivities and therefore we have not quoted it (e.g., for BP3 d_e/e is below 5×10^{-33} cm).

Conclusions. To conclude, we have shown that successful leptogenesis in the scotogenic model is possible for masses as low as 1.3 (2.5) TeV for thermal (zero) initial abundance of the lightest sterile neutrino, and for Yukawa couplings large enough to induce CLFV processes in current and planned experiments. The regions of parameter space we

have uncovered—which are compatible with neutrino masses and scalar DM—do not involve quasidegenerate neutrinos. Instead, washouts are suppressed exponentially as the result of the sizable mass of the inert Higgs and the action of some spectator processes. The treatment presented here for this suppression mechanism, which involves a massive particle different from its antiparticle, may apply to baryogenesis in other extensions of the SM.

Some directions for further development include a more exhaustive and precise exploration of the parameter space, considering, among other effects, the running of Yukawa couplings from the scale of leptogenesis to the low energy scales relevant for neutrino oscillation experiments.

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