Examining leptogenesis with lepton flavor violation
and the dark matter abundance

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\textbf{Abstract:} Within a supersymmetric (SUSY) type-I seesaw framework with flavor-blind universal boundary conditions, we study the consequences of requiring that the observed baryon asymmetry of the Universe be explained by either thermal or non-thermal leptogenesis. In the former case, we find that the parameter space is very constrained. In the bulk and stop-coannihilation regions of mSUGRA parameter space (that are consistent with the measured dark matter abundance), lepton flavor-violating (LFV) processes are accessible at MEG and future experiments. However, the very high reheat temperature of the Universe needed after inflation (of about $10^{12}$ GeV) leads to a severe gravitino problem, which disfavors either thermal leptogenesis or neutralino dark matter. Non-thermal leptogenesis in the preheating phase from SUSY flat directions relaxes the gravitino problem by lowering the required reheat temperature. The baryon asymmetry can then be explained while preserving neutralino dark matter, and for the bulk or stop-coannihilation regions LFV processes should be observed in current or future experiments.
1. Introduction

Supersymmetry is perhaps the leading possibility for physics beyond the Standard Model. One of its nice features is that it contains natural candidates for the observed dark matter in the Universe. Within the supersymmetric model with minimal particle content (MSSM), it is customary to assume flavor-blind boundary conditions at the Grand Unified Theory (GUT) scale, in which case the model is referred to as mSUGRA or constrained MSSM (CMSSM). If the dark matter particle is the lightest neutralino, the parameter space of mSUGRA is very tightly constrained by the precisely determined dark matter abundance in the Universe [1].

With exact R-parity conservation, neutrinos are massless in mSUGRA. However, there is now overwhelming evidence that neutrinos have mass and mix; for a review see Ref. [2]. The simplest explanation for small neutrino masses is perhaps the type-I seesaw mechanism [3]. It is therefore natural to extend mSUGRA to allow for a seesaw mechanism and thus for small neutrino masses. An mSUGRA-seesaw with SO(10)-inspired boundary conditions was recently studied in [4], where it was found that neutrinos, with their Yukawa couplings contributing to the running of various parameters, such as the slepton mass matrices and the trilinear couplings, substantially modify the parameter space allowed by dark matter. Lepton flavor violation (LFV) was then studied within the same framework, and the LFV rates were shown to potentially differ from existing estimates by up to two orders of magnitude [5].

In this paper we add yet another constraint to the mSUGRA-seesaw+dark-matter scenario, namely that the baryon asymmetry of the Universe be explained by either thermal or non-thermal leptogenesis.

It is well-known that in SO(10)-inspired scenarios where the type-I seesaw mechanism provides the dominant contribution to neutrino masses, thermal leptogenesis [6] typically fails to explain the observed baryon asymmetry of the Universe. The reason is that the lightest right-handed (RH) neutrino is generally too light to generate enough asymmetry [7]. Including flavor effects [8, 9], the situation improves since the next-to-lightest RH neutrinos (not accounted for in Ref. [7]), can generate a large asymmetry [10]. Nevertheless, the scenario remains tightly constrained [11, 12], which perfectly suits our purpose: If thermal leptogenesis is successful only in a very restricted part of the parameter space, then definite predictions for LFV rates at a given point in the mSUGRA parameter space are possible.

Thermal leptogenesis with hierarchical RH neutrinos requires the reheat temperature after inflation to be above $10^9$ GeV [13–15]. In mSUGRA this poses a problem because of the overproduction of gravitinos [16]. This tension is partially alleviated if the gravitino is heavier than 30 TeV, as the reheat temperature is then allowed to be as high as $10^{10}$ GeV [17]. As we show, such a reheat temperature is not high enough to allow for thermal leptogenesis from the next-to-lightest RH neutrino decays. The consequence is that thermal leptogenesis in our mSUGRA SO(10)-inspired framework is inconsistent with neutralino dark matter (or even gravitino dark matter).

The alternative possibility of non-thermal leptogenesis, either at reheating from inflaton decay [18, 19], or at preheating [20] allows for lower reheat temperatures than thermal
leptogenesis. We employ the mechanism of instant preheating [21] from SUSY flat directions, as presented in [22]. We show that this mechanism is able to successfully explain the baryon asymmetry of the Universe, while maintaining the viability of the neutralino or the gravitino as dark matter candidates. Our predictions for LFV rates turn out to be close to the current bounds for mSUGRA points in the bulk region.

It is worth mentioning that a mixed type-I + type-II [23] seesaw mechanism can be naturally obtained within SO(10), and leptogenesis becomes much easier [11, 24]. However, for the sake of minimalism and the predictiveness, we limit ourselves to a dominant type-I case only.

In Section 2 we describe the framework in which we work. In Section 3 we review thermal leptogenesis, introducing all the necessary tools for our computation. We also show the numerical results for the predicted LFV rates, and comment on the gravitino problem. In Section 4 we perform the same analysis with non-thermal leptogenesis at preheating. In Section 5 we summarize our findings and conclude.

2. SUSY-seesaw and SO(10) GUTs

We consider the following superpotential for the MSSM augmented by singlet right-handed neutrinos \( \tilde{N}_i^c \):

\[
\hat{f} = \hat{f}_{\text{MSSM}} + (f_\nu)_{\alpha j} \epsilon_{a b} \hat{L}_a^\alpha \hat{H}_u^b \tilde{N}_i^c + \frac{1}{2} (M_N)_{i j} \tilde{N}_i^c \tilde{N}_j^c ,
\]

where \( \alpha \) is the lepton flavor index, \( i, j \) are generation indices, \( a, b \) are \( SU(2)_L \) doublet indices, \( \epsilon_{a b} \) is the totally antisymmetric tensor with \( \epsilon_{12} = 1 \), and the superscript \( c \) denotes charge conjugation. Here, \( \hat{f}_{\text{MSSM}} \) is the MSSM superpotential, \( \hat{L} \) and \( \hat{H}_u \) are, respectively, the lepton doublet and up-Higgs superfields, and \( M_N \) is the Majorana mass matrix for the (heavy) right-handed neutrinos. At energy scales above \( M_N \), the light neutrino mass matrix is given by the type-I seesaw formula [3],

\[
\mathcal{M}_\nu = -f_\nu M_N^{-1} f_\nu^T v_u^2 ,
\]

where \( v_u \) is the vacuum expectation value (VEV) of the neutral component \( h_u^0 \) of the up-type Higgs doublet \( H_u \). We denote the eigenvalues of the light neutrino mass matrix by \( m_\nu i \), \( i = 1, 2, 3 \). In the limit in which all RH neutrinos are decoupled, the light neutrino mass matrix is \( \mathcal{M}_\nu = -\kappa v_u^2 \); \( \kappa \) is the coupling matrix of the dimension-5 effective operator generated by RH neutrinos, which is determined by matching conditions at the RH neutrino decoupling thresholds. The matrix \( \mathcal{M}_\nu \) is diagonalized (in the basis where charged leptons are diagonal) by the MNS matrix that can be parameterized by the mixing angles \( \theta_{12}, \theta_{23}, \) and \( \theta_{13} \), the Dirac phase \( \delta \), and two Majorana phases \( \phi_1 \) and \( \phi_2 \) (see Ref. [5] for our convention).

Inspired by SO(10) GUTs, we introduce the vector \( R_{\nu i} = (R_1, R_2, R_3) \), with strictly positive entries, which relates the diagonal up-type quark Yukawa couplings \( f_\nu^{\text{diag}} \) to the diagonal neutrino Yukawa couplings \( f_\nu^{\text{diag}} \) at the GUT scale:

\[
(f_\nu^{\text{diag}})^{ij} = R_i \left( f_\nu^{\text{diag}} \right)_j .
\]
Since we will always assume \( R_i < O(5) \), a highly hierarchical pattern \( f_{\nu_3}^{\text{diag}} \gg f_{\nu_2}^{\text{diag}} \gg f_{\nu_1}^{\text{diag}} \) is obtained. Note that the \( R_i \)'s are in general all different, and in the minimal \( SO(10) \) scenario, the range is typically \( 1 \leq R_i \leq 3 \). However, higher values \( R_i \gtrsim 5 \) can be easily achieved with non-renormalizable operators. The condition in Eq. (2.3) corresponds to an extension of the “small mixing” scenario presented in [5]. We will not consider the “large mixing” case here, since large regions of the parameter space are already excluded by existing bounds on \( \tau \rightarrow \mu \gamma \) [5].

Equation (2.3) implies that the RH neutrinos have a very strong hierarchy. To see this, assume tribimaximal mixing for the light neutrinos, and neglect the small CKM-type mixing in \( f_\nu \). For a normal hierarchy of light neutrinos, \( m_{\nu_1} \ll m_{\nu_2} \ll m_{\nu_3} \), one obtains [5]:

\[
M_{N_1} \simeq \frac{3m_{\nu_2}^2}{m_{\nu_2}} R_1^2, \quad M_{N_2} \simeq \frac{2m_{\nu_3}^2}{m_{\nu_1}} R_2^2, \quad M_{N_3} \simeq \frac{m_{\nu_1}^2}{6m_{\nu_3}} R_3^2, \tag{2.4}
\]

whereas for the inverted mass hierarchy (\( m_{\nu_1} \approx m_{\nu_2} \gg m_{\nu_3} \)),

\[
M_{N_1} \simeq \frac{3m_{\nu_2}^2}{m_{\nu_2}} R_1^2, \quad M_{N_2} \simeq \frac{2m_{\nu_3}^2}{3m_{\nu_1}} R_2^2, \quad M_{N_3} \simeq \frac{m_{\nu_1}^2}{2m_{\nu_3}} R_3^2. \tag{2.5}
\]

From the above scaling behavior, we see that a quasi-degenerate spectrum (\( m_{\nu_1} \approx m_{\nu_2} \approx m_{\nu_3} \)) would require the lightest Majorana mass to be in the \( 10^2 \text{–} 10^3 \) GeV range with significant L-R mixing in the sneutrino sector, which we disregard because it would substantially complicate the sneutrino mass spectrum and phenomenology. Moreover, since the next-to-lightest RH neutrino is also lighter than in the case of the normal hierarchy, successful thermal leptogenesis is rendered more difficult. The inverse hierarchical case would require the heaviest Majorana mass to be of order \( 10^{17} \) GeV, which is well above the GUT scale. This type of spectrum also suffers from instabilities under very small changes to \( M_N \) and RGE evolution [25]. For all these reasons we choose to focus on the normal hierarchy of light neutrinos, \( m_{\nu_1} \ll m_{\nu_2} \ll m_{\nu_3} \).

3. Thermal leptogenesis

Thermal leptogenesis is one of the most popular mechanisms to explain the observed baryon asymmetry of the Universe [6]. The crucial parameters for leptogenesis are the CP asymmetries \( \varepsilon_{i\alpha} \) and the washout parameters \( K_{i\alpha} \). The CP asymmetry from the decay of the heavy (s)neutrino \( N_i \) (\( \tilde{N}_i \)) into a (s)lepton of flavor \( \alpha \) is in full generality given by [26]

\[
\varepsilon_{i\alpha} = \frac{1}{8\pi(f_\nu^* f_\nu)_{ii}} \sum_{j \neq i} \left\{ \text{Im} \left[ (f_\nu^*)_{\alpha i} (f_\nu)_{\alpha j} (f_\nu^* f_\nu)_{ij} \right] g(x_j/x_i) + \frac{2}{(x_j/x_i - 1)} \text{Im} \left[ (f_\nu^*)_{\alpha i} (f_\nu)_{\alpha j} (f_\nu^* f_\nu)_{ji} \right] \right\}, \tag{3.1}
\]

where \( x_i \equiv M_{\tilde{N}_i}^2/M_{N_i}^2 \) and

\[
g(x) = \sqrt{x} \left[ \frac{2}{x - 1} + \ln \left( \frac{1 + x}{x} \right) \right]. \tag{3.2}
\]
In the SUSY limit, the decay width is
\[
\Gamma(N_i \to \ell_\alpha H_u) + \Gamma(N_i \to \bar{\ell}_\alpha H_u^c) = \Gamma(\tilde{N}_i^* \to \ell_\alpha H_u) = \Gamma(\tilde{N}_i \to \bar{\ell}_\alpha H_u) = \frac{|(f_\nu)_{\alpha i}|^2}{8\pi} M_{N_i}.
\]

We can then define
\[
K_{\alpha i} \equiv \frac{\Gamma(N_i \to \ell_\alpha H_u) + \Gamma(N_i \to \bar{\ell}_\alpha H_u^c)}{H(T = M_{N_i})} = \frac{v_u^2}{m_* M_{N_i}} |(f_\nu)_{\alpha i}|^2, \quad (3.3)
\]
where \(m_* \simeq (1.56 \times 10^{-3} \text{ eV}) \sin^2 \beta\).

The baryon asymmetry is obtained by solving a set of coupled Boltzmann equations as given for instance in Ref. [11]. However, we use convenient semi-analytical expressions for the final baryon asymmetry. The quantity that describes how efficiently the asymmetry is produced is the efficiency factor \(\kappa\), which is a function of \(K_{\alpha i}\). For an initial thermal abundance of RH (s)neutrinos, \(\kappa\) is given by [15],
\[
\kappa(K_{\alpha i}) \equiv \frac{1}{K_{\alpha i} z_B(2K_{\alpha i})} \left[ 1 - \exp \left( -\frac{2K_{\alpha i} z_B(2K_{\alpha i})}{2} \right) \right], \quad (3.4)
\]
where
\[
z_B(K) \simeq 2 + 4 K^{0.13} \exp \left( -\frac{2.5}{K} \right). \quad (3.5)
\]

With a vanishing initial abundance of RH (s)neutrinos, a different result ensues. A fit valid both in the weak washout (\(K_{\alpha i} < 3\)) and in the strong washout regime (\(K_{\alpha i} > 3\)), was obtained in [27]:
\[
\kappa(K_{\alpha i}) \simeq \left( \frac{2.6}{K_{\alpha i}} + \left( \frac{K_{\alpha i}}{0.06} \right)^{1.16} \right)^{-1}. \quad (3.6)
\]

For the hierarchical mass spectrum of the RH (s)neutrinos \(M_{N_1} \ll M_{N_2} \ll M_{N_3}\), the asymmetry production from each RH (s)neutrino can be considered separately, and eventually summed to obtain the final asymmetry. There are three mass ranges that need to be considered: the three-flavor regime for \(M_{N_i} < (1 + \tan^2 \beta) \times 10^9 \text{ GeV}\), the two-flavor (\(e + \mu\) and \(\tau\)) regime for \((1 + \tan^2 \beta) \times 10^9 \text{ GeV} < M_{N_i} < (1 + \tan^2 \beta) \times 10^{12} \text{ GeV}\) and the unflavored regime for \(M_{N_i} > (1 + \tan^2 \beta) \times 10^{12} \text{ GeV}\) [8, 9].

It is well-known that with \(SO(10)\)-inspired mass relations such as in Eq. (2.3), the lightest RH (s)neutrino \(N_1 (\tilde{N}_1)\) typically fails to produce enough asymmetry because \(M_{N_1}\) is usually predicted to be much smaller than the Davidson-Ibarra bound [13] of about \(10^9 \text{ GeV}\) for successful leptogenesis. Note however that tiny regions in the parameter space exist where RH neutrino masses are quasi-degenerate, \(M_{N_1} \simeq M_{N_2} \simeq M_{N_3}\), in which case the \(CP\) asymmetry can be dramatically enhanced [28] and leptogenesis is possible [7]. We do not entertain this possibility any further.

\footnote{For the SM case, the replacement \(2 K_{\alpha i} \to K_{\alpha i}\) must be made because of the fewer decay modes for each heavy particle.}
We consider the crucial contribution to leptogenesis to arise from the next-to-lightest RH (s)neutrino, $N_2$ ($\tilde{N}_2$) [10]. The asymmetry is typically produced in the two-flavor regime, but it is necessary to include the potential washout from $N_1$ ($\tilde{N}_1$), which occurs in the three-flavor regime. We then have [12, 29]

$$\eta_{B,2} \simeq 0.96 \times 10^{-2} \left[ \tilde{\varepsilon}_{2e} \kappa(K_{2e} + K_{2\mu}) \exp \left( -\frac{3\pi}{4} K_{1e} \right) + \tilde{\varepsilon}_{2\mu} \kappa(K_{2e} + K_{2\mu}) \exp \left( -\frac{3\pi}{4} K_{1\mu} \right) 
+ \tilde{\varepsilon}_{2\tau} \kappa(K_{2\tau}) \exp \left( -\frac{3\pi}{4} K_{1\tau} \right) \right].$$

(3.7)

If the asymmetry from $N_2$ ($\tilde{N}_2$) is produced in the three-flavor regime,

$$\eta_{B,2} \simeq 0.96 \times 10^{-2} \left[ \tilde{\varepsilon}_{2e} \kappa(K_{2e}) \exp \left( -\frac{3\pi}{4} K_{1e} \right) + \tilde{\varepsilon}_{2\mu} \kappa(K_{2\mu}) \exp \left( -\frac{3\pi}{4} K_{1\mu} \right) 
+ \tilde{\varepsilon}_{2\tau} \kappa(K_{2\tau}) \exp \left( -\frac{3\pi}{4} K_{1\tau} \right) \right].$$

(3.8)

Since the asymmetry production from the heaviest RH (s)neutrino, $N_3$ ($\tilde{N}_3$), occurs typically in the unflavored regime, the $CP$ asymmetry is suppressed by $(M_{N_1,2}/M_{N_3})^2$, and we neglect it.

The total asymmetry is given by the sum of the three RH neutrino contributions, of which just one is relevant, and thus

$$\eta_B \simeq \eta_{B,2},$$

(3.9)

to be compared with the measured value [1]

$$\eta_B^{\text{CMB}} = (6.2 \pm 0.15) \times 10^{-10}.$$  

(3.10)

As we shall see in the next subsection, the final baryon asymmetry produced through leptogenesis will be typically dependent on the initial abundance of RH (s)neutrinos; Eq. (3.4) is valid for an initial thermal abundance of $N_2$ ($\tilde{N}_2$) whereas Eq. (3.6) is valid for a vanishing one. Note that within our $SO(10)$-inspired scenario, a thermal $N_2$ ($\tilde{N}_2$)-abundance is very easily obtained if the $Z'$ of $U(1)_{B-L}$, which is naturally present if $SO(10)$ breaks to the left-right model, is heavier than $N_2$ ($\tilde{N}_2$) by two orders of magnitude (and $M_{Z'} \lesssim 10 T_R$, with $T_R$ the reheat temperature) [30]. We then have that the interaction $\overline{N_2} \gamma^\mu N_2 Z'_\mu$ efficiently brings the next-to-lightest RH (s)neutrino into equilibrium without interfering with the production mechanism from $N_2$ ($\tilde{N}_2$) decays.

### 3.1 Results

As stated in Section 2, we present results only for the normal hierarchy of light neutrinos, $m_{\nu_1} < m_{\nu_2} < m_{\nu_3}$. More precisely, we require that $m_{\nu_1} < m_{\text{sol}} = 0.009$ eV.

We use ISAJET-M to produce the neutrino and SUSY spectra [5]. The program implements RGE evolution in the MSSM with a type-I seesaw in full matrix form at the 2-loop level. All sparticle masses are computed with complete 1-loop corrections and gauge and Yukawa coupling evolution include multiple sparticle threshold effects. The decoupling
of RH neutrinos is performed at multiple scales equal to their own running masses, and effects of $f_\nu$ on the MSSM parameters are included. It is known that large $f_\nu$ entries can significantly affect the RGE evolution with concomitant effects on the MSSM spectrum and the neutralino DM rates \[4, 5, 31\].

In the neutrino sector, we use a top-down approach with $f_\nu$ and $M_N$ input at $M_{\text{GUT}}$ and the neutrino mass matrix $M_\nu$ obtained by RGE evolution. We fix $f_\nu(M_{\text{GUT}})$ using the $SO(10)$-inspired relation (2.3) and adjust $M_N$ to produce a viable spectrum of light neutrinos. Since leptogenesis is sensitive to details of the MNS matrix, we scan over 5 parameters: $\phi_1, \phi_2, \delta, \theta_{13}$ and $m_{\nu_1}$.

In Fig. 1, we show the results for the baryon asymmetry generated by thermal leptogenesis for $m_{\nu_1} = 0.005$ eV and $R_\nu = (1, 5, 1)$ with $\phi_1, \phi_2, \delta$ and $\theta_{13}$ varied. Note that the figure displays weak scale values for all the parameters. From the bottom-right panel, we see that there are only a few points above the 2$\sigma$ lower bound on the observed baryon abundance [cf. Eq. (3.10)], i.e., $\eta_B > 5.9 \times 10^{-10}$. This clearly shows that the parameter space that yields successful leptogenesis is quite restricted.

As noted in Ref. [12], the baryon asymmetry is essentially independent of $R_1$, and only mildly dependent on $R_3$. (At the end of the section we discuss the role of $R_3$ on the predictions for the LFV rates.) On the other hand, $R_2$ is crucial for leptogenesis, in that it fixes the next-to-lightest RH neutrino mass scale [cf. Eq. (2.4)], which itself sets the size of the CP asymmetry [cf. Eq. (3.1)]. We confirm the finding of Ref. [12] that leptogenesis is only possible for $R_2 > 3$, and set $R_2 = 5$ in our calculations.

To obtain more points with sufficiently large values of the baryon asymmetry, we now focus on a restricted parameter space. In Fig. 2, we show results for $\eta_B$ in the parameter space, $2 < \phi_1 < 4, -1 < \phi_2 < 1, 1 < \delta < 5$ and $0 < \theta_{13} < 0.2$, with $m_{\nu_1}$ varied. We obtain many allowed points, most of which have in common that the asymmetry is produced in the e flavor, and where $K_{2e} + K_{2\mu} \sim 1/2$, in which case the efficiency factor is close to maximal, and depends mildly on the initial conditions. With a vanishing initial RH neutrino abundance the efficiency factor would be lower by a factor of 2–3 for these points, so that the final asymmetry would be slightly lower than that observed. Note that although the $CP$ asymmetry $\varepsilon_{2\tau}$ is typically the largest one (and so is $K_{2\tau}$), the asymmetry in the tau flavor typically suffers from a large washout from $K_{2\tau}$; the washout from $N_1(\tilde{N}_1)$, although set by a relatively small $K_{1\tau} \lesssim 20$, can also have a large impact due to its exponential effect [see Eq. (3.7)].

From Fig. 2 we find a lower bound on the lightest neutrino mass and on $\theta_{13}$:

\begin{equation}
    m_{\nu_1} \gtrsim 0.004 \text{ eV}, \quad \theta_{13} \gtrsim 0.04. \tag{3.11}
\end{equation}

Our lower bound on the lightest neutrino mass is slightly more restrictive than found in Ref. [12], where a non-SUSY framework and a vanishing RH neutrino abundance were considered. On the other hand, our lower bound on $\theta_{13}$ agrees well with that of Ref. [12] in the mass region below $m_1 = 0.009$ eV. Note that the Daya Bay [32] and Double Chooz [33] reactor experiments are sensitive to $\theta_{13}$ for $\theta_{13} \gtrsim 0.05$.

As mentioned above, with a vanishing initial $N_2(\tilde{N}_2)$-abundance it is very difficult to obtain a baryon asymmetry in the allowed range. At this point, it is worth commenting on
Figure 1: Full parameter space scan for thermal leptogenesis (thermal initial $N_2$-abundance). The color code is evident from the bottom-right panel.

the size of the theoretical errors in the computation of the baryon asymmetry. First, we have neglected spectator processes [34], including flavor mixing in the so-called $C$ matrix [8], which can reduce the final asymmetry by up to 30% [29]. We also neglected quantum
Figure 2: Baryon asymmetry from thermal leptogenesis (thermal initial $N_2$-abundance) vs. $\theta_{13}$ for different values of the lightest neutrino mass $m_{\nu_1}$. The horizontal line marks the 2$\sigma$ lower bound, $\eta_B > 5.9 \times 10^{-10}$.

statistical factors and assumed that kinetic equilibrium holds, an approximation that is very good in the strong washout regime ($K_{1\alpha} \gg 1$), but which can make a 50% difference in the weak washout regime ($K_{1\alpha} \ll 1$) [35]. In our study, we obtained the largest asymmetries when $K_{1\alpha} \sim 1$, in which case the uncertainty is less than 10%. Finally, we did not solve the full quantum Boltzmann equations based on the Keldysh-Schwinger non-equilibrium formalism; for recent related work see Ref. [36]. Note that in this formalism, thermal corrections are automatically taken into account. According to Ref. [37] an enhancement of the $CP$ asymmetry parameter by a factor of a few is possible.

It is difficult to estimate the cumulative effect of all these theoretical uncertainties on our results. Nevertheless, a 50% uncertainty in the final asymmetry seems to be a fair assessment and we conclude that a vanishing initial RH neutrino abundance in this framework cannot be excluded.

Next, we extract from Fig. 2 the values of the RH neutrino masses $M_{N_2}$ required for successful thermal leptogenesis. The results are shown in Fig. 3. These clearly point to a very high reheat temperature, $T_R \sim M_{N_2} \sim 10^{12}$ GeV, which poses a problem as we explain in the next subsection.

We now turn to LFV. It has been known for a long time that the supersymmetric seesaw potentially leads to large LFV rates, due largely to slepton contributions at the loop level [38]. It is then a quantitative question to know if the points compatible with
Figure 3: Values of $M_{N_2}$ for which thermal leptogenesis is successful for different values of the lightest neutrino mass $m_{\nu_1}$.

thermal leptogenesis in our framework lead to observable rates in future experiments. We show in Table 1 the current bounds and projected sensitivities for LFV.

So far our results have been essentially independent of the region of mSUGRA parameter space compatible with the dark matter abundance since leptogenesis occurs at very high energy scales. However, for the LFV rates, it is of crucial importance. We focus on the bulk region, which is the most optimistic region for the detection of LFV. This region is characterized by the following parameters: $m_0 = 80$ GeV, $m_{1/2} = 170$ GeV, $A_0 = -250$ GeV and $\tan \beta = 10$. We also checked that points in the stop-coannihilation region ($m_0 = 150$ GeV, $m_{1/2} = 300$ GeV, $A_0 = -1095$ GeV and $\tan \beta = 5$) yield rates that are no more than a factor of two different than in the bulk region [5]. This is not surprising given the fact that LFV rates are maximized when sfermions and gauginos are light, and when the mass hierarchy is mild between them.

Table 1: Present bounds and projected sensitivities for LFV processes.

| Process            | Present       | Future       |
|--------------------|---------------|--------------|
| $\text{BR}(\mu \to e\gamma)$ | $1.2 \times 10^{-11}$ [39] | $10^{-13}$ [45] |
| $\text{BR}(\tau \to \mu\gamma)$ | $4.5 \times 10^{-8}$ [40] | $10^{-9}$ [46] |
| $\text{BR}(\tau \to e\gamma)$ | $3.3 \times 10^{-8}$ [41] | $10^{-9}$ [46] |
| $\text{BR}(\mu \to eee)$ | $1.0 \times 10^{-12}$ [42] | $10^{-14}$ [47] |
| $\text{BR}(\tau \to \mu\mu\mu)$ | $3.2 \times 10^{-8}$ [43] | $10^{-9}$ [46] |
| $\text{BR}(\tau \to eee)$ | $3.6 \times 10^{-8}$ [43] | $10^{-9}$ [46] |
| $\text{CR}(\mu \text{Ti} \to e \text{Ti})$ | $4.3 \times 10^{-12}$ [44] | $10^{-18}$ [48] |
| $\text{CR}(\mu \text{Al} \to e \text{Al})$ | - | $10^{-16}$ [49] |
**Bulk region,** $R_{\nu_1^u}=(1,5,1), m_{\nu_1}=0.005$ eV, $\eta_B \geq 5.9 \times 10^{10}$

- **thermal**
- **preheating** $T_R = 10^8$ GeV

Figure 4: LFV rates for points with an asymmetry above $5.9 \times 10^{-10}$ within thermal leptogenesis (red dots) and non-thermal leptogenesis at preheating (blue stars).

The results for LFV rates in the bulk region are presented in Fig. 4. We find that the allowed points lead to predictions for LFV rates below the present exclusion limits,
and, except for $\tau \to \mu \mu \mu$, within reach of current and future experiments. In particular, MEG should see a $\mu \to e\gamma$ signal if leptogenesis is realized in our framework. It should be noted that LFV rates are very weakly dependent on the value of the lightest neutrino mass $m_{\nu_1}$, and we checked that the rates are essentially unchanged in the range $0.004$ eV $< m_{\nu_1} < 0.009$ eV.

We end this section by commenting on the dependence of LFV rates on $R_3$. Increasing $R_3$ leads to a larger $M_{N_3}$ [see Eq. (2.4)], which in turn leads to larger Yukawa couplings $(f_\nu)_{\alpha 3}$ in order to keep the neutrino mass matrix fixed. We expect LFV rates to increase as we increase $R_3$ since the third generation Yukawa contributes dominantly to the rates. We explicitly checked that this is the case, and found the rates for $R_3 = 5$ to be about one order of magnitude larger than those shown in Fig. 4.

### 3.2 Reheat temperature and the gravitino problem

In order for leptogenesis to work in our framework, we need a very high reheat temperature ($\sim 10^{12}$ GeV) of the Universe after inflation (see Fig. 3). Therefore, the gravitino problem here is even more severe than in the conventional scenario, where typically a reheat temperature of about $10^{9}$ GeV is sufficient.

Let us first consider the case where the gravitino is not the lightest supersymmetric particle (LSP). The gravitino problem arises because the thermal production of gravitinos is unavoidable after inflation. The gravitino yield, $Y_{3/2}$, defined as the gravitino number density divided by the entropy density, is nearly proportional to the reheat temperature [17]:

$$ Y_{3/2} \simeq 2.3 \times 10^{-14} \times T_R^{(8)} \left[ 1 + 0.015 \log T_R^{(8)} - 0.0009 \log^2 T_R^{(8)} \right]$$

$$+ 1.5 \times 10^{-14} \times \left( \frac{m_{1/2}}{m_{3/2}} \right)^2 T_R^{(8)} \left( 1 - 0.037 \log T_R^{(8)} + 0.0009 \log^2 T_R^{(8)} \right), \quad (3.12)$$

where $T_R^{(8)} \equiv T_R/(10^8$ GeV), $m_{3/2}$ is the gravitino mass, and $m_{1/2}$ is the unified gaugino mass at the GUT scale. Note also that increasing the gravitino mass lowers its abundance until the first term in Eq. (3.12) dominates, at which point the abundance saturates.

Gravitinos being only gravitationally coupled, for masses below 30 TeV they typically decay during or after Big Bang Nucleosynthesis (BBN), hence spoiling the agreement between observations and theory. Requiring that BBN is successful leads to stringent constraints on the reheat temperature as a function of the gravitino mass (see [17] and references therein): $T_R < 10^6$ GeV if the gravitino mass is lower than about 10 TeV.

For $m_{3/2} \gtrsim 30$ TeV, the bound relaxes to $T_R \lesssim 10^{10}$ GeV so that that the gravitino decays into the LSP (assumed to be the lightest neutralino) yielding a non-thermal contribution that saturates the observed dark matter density. This upper bound is still at odds with thermal leptogenesis in our framework. Therefore, we conclude that thermal leptogenesis is somewhat incompatible with neutralino dark matter.

Suppose we abandon neutralino dark matter. Consistent cosmology first requires that the gravitino be heavier than 30 TeV so that it decays to the LSP before the onset of BBN. In turn, the LSP may decay to a hidden sector [50]. The lightest hidden sector
particle $X$ needs to be much lighter than the LSP so that its energy density is diluted before matter-radiation equality:

$$\Omega_X h^2 \simeq 2.8 \times 10^5 \times Y_{3/2} \left( \frac{m_X}{1 \text{ MeV}} \right).$$  

Using Eq. (3.12) for a very heavy gravitino $m_{3/2} \gg m_{1/2}$, we find

$$\Omega_X h^2 \simeq 6 \times 10^{-5} \left( \frac{T_R}{10^{12} \text{ GeV}} \right) \left( \frac{m_X}{1 \text{ MeV}} \right).$$  

The hidden sector particle could constitute part of the dark matter, but being warm, its abundance cannot exceed 5% of the total dark matter abundance [51].

An alternative way to circumvent the gravitino problem is if the gravitino is the (visible) LSP, with a mass of about 200–300 GeV, and itself decays into a much lighter hidden sector particle long before matter-radiation equality. In this case, the NLSP must decay before BBN, which can be easily achieved with hidden sector dynamics [50]. Note that the gravitino cannot be the dark matter particle.

Independently of $m_{3/2}$, we find that dark matter must be explained by an external mechanism and particle, like the axion.

### 4. Non-thermal leptogenesis at preheating

As explained in the previous section, within our mSUGRA-seesaw framework with $SO(10)$-inspired boundary conditions, standard thermal leptogenesis is subject to a severe gravitino problem.

Non-thermal leptogenesis allows for a low reheat temperature so that neutralino dark matter is viable. Non-thermal RH neutrino production can be obtained either during the preheating stage [20], or from inflaton decays [18, 19].

Here we follow the approach of Ref. [22] which does not require any additional ingredient to our framework such as a large coupling of the RH neutrinos to the inflaton. The idea is to use the presence of flat directions in the scalar potential (for a review see [52]) to enable instant preheating [21], which is a very efficient way of producing very heavy states. We briefly review the mechanism.

$F$- and $D$-term flat directions can be lifted by soft supersymmetry (SUSY) breaking terms in our vacuum, with non-renormalizable terms in the superpotential, or with finite density terms in the potential [53]. If the soft SUSY-breaking term has a positive sign and the finite density term (proportional to the Hubble rate squared $H^2$), contributes negatively, the field $\phi$ along the flat direction initially acquires a large vacuum expectation value (VEV) denoted by $\phi_0$. After inflation ends, the inflaton starts oscillating at the minimum of its potential, while the Hubble rate falls. Once $H \sim \tilde{m}/3$, where $\tilde{m}$ is a soft SUSY-breaking mass term, the flat direction starts moving down towards the true minimum at $\phi = 0$. Following [22], we assume that the condensate involves the third generation quark $u_3$, and focus on the production of the up-type scalar Higgs $H_u$ relevant for leptogenesis. The condensate couples to the Higgs through the term $f_t|\phi|^2|H_u|^2$. If the
condensate passes through the origin (or sufficiently close to it), Higgses will be produced when adiabaticity is violated [21], i.e. \( \dot{m}_{H_u}/m_{H_u}^2 \gtrsim 1 \).

The condensate continues its motion upwards after it has passed through the origin, and the up-type Higgs effective mass gradually increases proportionally to \( f_t|\phi| \). When the Higgs effective mass becomes larger than the RH neutrino mass, it promptly decays to RH neutrinos. The heaviest particles that can be produced through this mechanism have a mass [22]

\[
M^{\text{max}} \simeq 4 \times 10^{12} \text{ GeV} \left( \frac{|\phi_0|}{M_{\text{Pl}}} \right)^{1/2} \left( \frac{\tilde{m}}{100 \text{ GeV}} \right)^{1/2},
\]

(4.1)

which is large enough for our purposes. Accounting for the fact that reheating occurs after leptogenesis, a large dilution factor must be included. Assuming that all up-type Higgses decay into RH neutrinos, the final result for the baryon asymmetry is given by [22]

\[
\eta_{B,i} \simeq 7 \times 10^{-5} \sum_{\alpha} \xi_{i\alpha} \left( \frac{T_R}{10^8 \text{ GeV}} \right) \left( \frac{|\phi_0|}{M_{\text{Pl}}} \right)^{3/2} \left( \frac{100 \text{ GeV}}{\tilde{m}} \right)^{1/2},
\]

(4.2)

which means that for the canonical choice of parameters the efficiency factor is about \( 7 \times 10^{-3} \).

For a highly hierarchical RH neutrino mass spectrum, and small \( M_{N_1} \), we need to again consider the production of asymmetry from \( N_2 \). Specifically, we need the up-type Higgs to promptly decay into \( N_2 \) rather than \( N_1 \), which is possible if [22]

\[
M_{N_2} \lesssim \left( \frac{8\pi f_t \tilde{m}|\phi_0|}{\sum_{\alpha} |(f_{\nu})_{\alpha 1}|^2} \right)^{1/2},
\]

(4.3)

The condition can be easily satisfied if \( |(f_{\nu})_{\alpha 1}| \ll 1 \).

Finally, we need to take into account the washout by \( N_1 \) inverse decays of the asymmetry produced by \( N_2 \). Using Eq. (4.2) we find

\[
\eta_{B,2} \simeq 7 \times 10^{-5} \sum_{\alpha} \xi_{2\alpha} \exp \left( -\frac{3\pi}{4} K_{1\alpha} \right),
\]

(4.4)

where we have taken \( T_R = 10^8 \text{ GeV}, \tilde{m} = 100 \text{ GeV} \) and \( \phi_0 = M_{\text{Pl}} \).

Before turning to the numerical results, we comment on another non-thermal scenario of leptogenesis, namely at reheating from inflaton decay [18, 19]. The final asymmetry obtained for an inflaton of mass \( 10^{13} \text{ GeV} \) is [18]

\[
\eta_{B,i} \simeq \frac{7}{2} \times 10^{-5} \sum_{\alpha} \xi_{i\alpha} \text{Br}(\phi \to N_iN_i) \left( \frac{T_R}{10^8 \text{ GeV}} \right),
\]

(4.5)

where \( \text{Br}(\phi \to N_iN_i) \) is the branching ratio of the inflaton decay channel \( \phi \to N_iN_i \). We find that even if the branching ratio to the next-to-lightest RH neutrino \( N_2 \) is unity, the maximum efficiency factor is a factor of 2 smaller than in the preheating case. All the results presented in the next subsection can then be trivially extended to the reheating case.
4.1 Results

Similarly to the previous section, we explore the parameter space for which Eq. (4.4) reproduces the observed baryon asymmetry; see Fig. 5. Note that the measured baryon asymmetry can be generated for more of the parameter space than with thermal leptogenesis. The reason is simply that the asymmetry depends on fewer parameters than in the thermal case. In particular, the asymmetry does not depend on $K_{2\alpha}$ as was the case for thermal leptogenesis [see Eqs. (4.4) and (3.7)]. Therefore, non-thermal leptogenesis is possible even for very large values of the washout parameter $K_{2\alpha}$, so long as $K_{1\alpha}$ remains small. In the non-thermal case, the $SO(10)$ mass relations in Eq. (2.3) and light neutrino masses lead us to a new part of the parameter space in which the Yukawa coupling $(f_\nu)_{\tau 2}$ (and therefore $K_{2\tau}$) is large, while $K_{1\tau}$ remains small. In this region, the lepton asymmetry is produced in the $\tau$ flavor because of the large value of the $CP$ asymmetry parameter $\varepsilon_{2\tau}$ and can be up to two orders of magnitude larger than the observed value. Due to the greater freedom in the choice of parameters, it is not surprising that neither $m_{\nu 1}$ nor $\theta_{13}$ are bounded from below.

For the points in Fig. 5, we show the corresponding LFV rates in Fig. 4. As for thermal leptogenesis, we find that large rates are predicted for the bulk and stop-coannihilation regions. In particular, MEG should see a positive signal if leptogenesis is the origin of the baryon asymmetry. Note that the rates remain essentially unchanged under variations of the lightest neutrino mass $m_{\nu 1}$ within the range of interest.

Finally, let us comment on the gravitino problem in this framework. Since the reheat temperature after inflation is required to be of order $10^7$–$10^8$ GeV, it is clearly not as severe as in the previous section. However, if dark matter is to be explained by the standard neutralino LSP, we still need a fairly heavy gravitino of about 10 TeV [17]. On the other hand, if the gravitino is the LSP, we must ensure that the NLSP decays before BBN, which can be easily achieved with small R-parity violation [54], or with decays into a light hidden sector [50].

5. Conclusions

We studied the implications of successful leptogenesis on the mSUGRA-seesaw parameter space with dark matter comprised of neutralinos. Guided by $SO(10)$-inspired mass relations, we were led to hierarchical Dirac mass eigenvalues in the neutrino sector.

We found that with thermal leptogenesis, a large enough baryon asymmetry is difficult to realize, and obtained lower bounds on both the lepton mixing angle $\theta_{13}$ and the lightest neutrino mass $m_{\nu 1}$. The LFV rates in the bulk and stop-coannihilation regions are large and observable at current experiments such as MEG. However, the high reheat temperature ($\sim 10^{12}$ GeV) implies a severe overproduction of gravitinos, rendering neutralino dark matter and thermal leptogenesis somewhat incompatible in our framework.

To relax the tension with gravitino overproduction, so as to not abandon neutralino dark matter, we explored the possibility of non-thermal leptogenesis in the preheating phase which relies on the mechanism of instant preheating from supersymmetric flat directions already present in our model. We found that an efficient (non-thermal) production of
Figure 5: Full parameter space scan for non-thermal leptogenesis at preheating. The color code is evident from the bottom-right panel.

RH neutrinos can be achieved, even for reheat temperatures as low as $10^7$ GeV. Note that non-thermal leptogenesis from inflaton decay would lead to very similar results. The parameter space for successful leptogenesis at preheating is less constrained than in the
thermal framework, and LFV rates in the bulk and stop-coannihilation regions remain large and observable at current and future experiments.

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