A new, direct link between the baryon asymmetry and neutrino masses

Michele Frigerio\textsuperscript{a}, Pierre Hosteins\textsuperscript{b}, Stéphane Lavignac\textsuperscript{a} and Andrea Romanino\textsuperscript{c}

\textsuperscript{a} Institut de Physique Théorique, CEA-Saclay, 91191 Gif-sur-Yvette Cedex, France
\textsuperscript{b} Department of Physics, University of Patras, GR-26500 Patras, Greece
\textsuperscript{c} International School for Advanced Studies (SISSA) and INFN, I-34013 Trieste, Italy

Abstract

We point out that, in a class of $SO(10)$ models with matter fields in 16 and 10 representations and type II realization of the seesaw mechanism, the light neutrino masses and the CP asymmetry needed for leptogenesis are controlled by one and the same set of couplings. The generated baryon asymmetry then directly depends on the low-energy neutrino parameters, with no unknown seesaw-scale flavour parameters involved; in particular, the necessary CP violation is provided by the CP-violating phases of the lepton mixing matrix. We compute the CP asymmetry in triplet decays for this scenario and show that it can lead to successful leptogenesis.

1 Introduction

The seesaw mechanism \cite{seesaw} nicely connects two pieces of observation that remain unexplained in the Standard Model: the smallness of neutrino masses and the fact that our universe contains almost no antimatter. Indeed, the same heavy states that are responsible for the generation of small Majorana masses for the Standard Model neutrinos in the seesaw mechanism can be the source of the observed baryon asymmetry through their decays. The lepton asymmetry generated in the out-of-equilibrium decays of the heavy states is then partially converted into a baryon asymmetry by the non-perturbative sphaleron processes \cite{sphaleron}. This mechanism is known as baryogenesis via leptogenesis \cite{leptogenesis}. The role of the heavy states is generally played by right-handed

\textsuperscript{1}Laboratoire de la Direction des Sciences de la Matière du Commissariat à l’Energie Atomique et Unité de Recherche Associée au CNRS (URA 2306).
neutrinos (type I seesaw mechanism) and/or by scalar $SU(2)_L$ triplets (type II seesaw mechanism). In spite of this intimate connection between neutrino masses and leptogenesis, there is little correlation, in general, between the generated baryon asymmetry and the neutrino parameters that can be measured at low energy. For instance, in the Standard Model augmented with a type I seesaw mechanism, the CP asymmetry in right-handed neutrino decays is not directly related to the CP-violating phases of the lepton mixing (PMNS) matrix. Even in the constrained framework of $SO(10)$ Grand Unified Theories (GUTs), there is generally enough freedom to spoil any relationship of this kind. Exceptions to this general statement are possible at the price of strong assumptions on the seesaw parameters.

In this paper, we present a new leptogenesis scenario in which the generated baryon asymmetry is directly related to the low-energy neutrino parameters, with no dependence on unknown seesaw-scale flavour parameters. In particular, the relevant CP asymmetry depends on the CP-violating phases of the PMNS matrix. This scenario is realized in $SO(10)$ models with Standard Model fermions split among $16$ and $10$ representations and type II realization of the seesaw mechanism.

The paper is organized as follows. In Section 2, we describe the class of $SO(10)$ models in which our leptogenesis scenario takes place. In Section 3, we present this scenario and compute the relevant CP asymmetry. We then discuss its dependence on the light neutrino parameters. In Section 4, we write the Boltzmann equations and estimate the efficiency factor that determines the final baryon asymmetry. We also comment on lepton flavour effects and on the contributions of other heavy states present in the model. Finally, we present our conclusions in Section 5.

2 The $SO(10)$ framework

In standard $SO(10)$ unification, all Standard Model fermions of a given generation reside in a $16$ representation, together with a right-handed neutrino. The charged fermion and Dirac neutrino mass matrices receive contributions from Yukawa couplings of the form $16_{i}^{\times} 16_{j}^{\times} H$ (where $H = 10, \overline{126}$ and/or $120$) as well as from non-renormalizable operators suppressed by powers of a scale $\Lambda \gg M_{GUT}$. Majorana masses for the right-handed neutrinos are generated either from $16_{i}^{\times} 16_{j}^{\times} \overline{126}$, or from the non-renormalizable operators $16_{i}^{\times} 16_{j}^{\times} 16_{i}^{\times} 16_{j}^{\times} / \Lambda$. This leads to small Majorana masses for the light neutrinos via the type I seesaw mechanism. In addition, light neutrino masses may also receive a type II seesaw contribution from the exchange of the electroweak triplet contained in the $\overline{126}$ representation. Leptogenesis in this framework has been studied by many authors, both in the type I and in the type I+II cases. In spite of the existence of constraining $SO(10)$ mass relations, the predictivity of these models for leptogenesis is generally limited, due in particular to the presence of physical high energy phases contributing to the CP asymmetry.

In this paper, we consider a different class of (supersymmetric) $SO(10)$ models, in which the Standard Model fermions are split among $16$ and $10$ matter multiplets. More precisely, we are interested in a subclass of these models in which neutrino masses arise from a pure type II seesaw mechanism. The advantage of these models is that,
as we are going to see, they lead to a more predictive leptogenesis. To be specific, let us consider the following superpotential:

\[
W = \frac{1}{2} y_{ij} \mathbf{16}_i \mathbf{16}_j \mathbf{10} + h_{ij} \mathbf{16}_i \mathbf{10}_j \mathbf{16} + \frac{1}{2} f_{ij} \mathbf{10}_i \mathbf{10}_j \mathbf{54} \\
+ \frac{1}{2} \sigma \mathbf{10} \mathbf{10} \mathbf{54} + \frac{1}{2} M_{54} \mathbf{54}^2 + \cdots ,
\]

where \( \mathbf{16}_i \) and \( \mathbf{10}_i \) \((i = 1, 2, 3)\) are matter representations, \( \mathbf{10}, \mathbf{16} \) and \( \mathbf{54} \) are Higgs representations, and we imposed a matter parity to restrict the form of the superpotential. The splitting of the matter representations \( \mathbf{16}_i \equiv (\mathbf{10}^{16}, \mathbf{5}^{16}, \mathbf{1}^{16})_i \) and \( \mathbf{10}_i \equiv (\mathbf{5}^{10}, \mathbf{5}^{10})_i \) is realized by the vev \( v_1^{16} \) of the \( SU(5) \)-singlet component of \( \mathbf{16} \), which gives large \( SU(5) \)-invariant masses \( M_i = h_i v_1^{16} \) to the \((5^{16}, 5^{10})\) pairs. In the following, we shall refer to the components of \( 5^{10}_i \) as heavy antilepton doublets \( L_i^{c} \) (with lepton number \( L = -1 \)) and heavy quark singlets \( D_i \) (with baryon number \( B = +1/3 \)), respectively:

\[
5^{10}_i \equiv (L_i^{c}, D_i).
\]

The Standard Model fermions are identified as:

\[
10^{16}_i = (Q, u_i^c, e_i^c) , \quad 5^{10}_i = (L, d_i^c),
\]

(3)

(The remaining \( \mathbf{16}_i \) component being the right-handed neutrino: \( 1^{16}_i = \nu_i^c \)), and their mass matrices are given by, at the renormalizable level:

\[
M_u = y v_1^{10} , \quad M_d = M_e^T = h v_1^{16} ,
\]

(4)

where \( v_1^{10} \) and \( v_1^{16} \) are electroweak-scale vevs associated with the \( SU(2)_L \) doublet components \( H_u^{10} \) and \( H_d^{16} \) of \( \mathbf{10} \) and \( \mathbf{16} \), respectively.\(^2\) As is well known, the GUT-scale mass relation \( M_d = M_e^T \) is not consistent with the measured values of the charged lepton and down quark masses, and must be corrected by non-renormalizable operators.

As for neutrino masses, they arise from a pure type II seesaw mechanism, in contrast with standard \( SO(10) \) unification where the type I seesaw contribution is always present. Indeed, Eqs. (1) and (3) imply that the Dirac mass matrix vanishes at the renormalizable level, while the \( \mathbf{54} \) Higgs multiplet contains a pair of electroweak triplets \((\Delta, \Sigma)\) with the requisite couplings \( \frac{1}{2} f_{ij} L_i L_j \Delta \) and \( \frac{1}{2} \sigma H_u^{10} H_d^{16} \Sigma \) to mediate a Majorana mass term for the Standard Model neutrinos:

\[
m_\nu = \frac{\sigma (v_1^{10})^2}{2M_\Delta} f ,
\]

(5)

with \( M_\Delta = M_{54} \). In this paper, we neglect possible subdominant type I contributions to Eq. (4) coming from non-renormalizable operators. Note that this realization of the type II seesaw mechanism is very different from standard \( SO(10) \) unification with a \( 126 + \overline{126} \) pair.

The above conclusions hold as long as the \( 5^{10}_i \)'s do not mix with the \( 5^{16}_i \)'s. This assumes in particular that the \( \mathbf{54} \) Higgs multiplet in Eq. (1) does not acquire a GUT-scale

\(^2\)We assume here that the light Higgs doublets \( H_u \) and \( H_d \) contain a significant component of \( H_u^{10} \) and \( H_d^{16} \), respectively. We checked that this naturally arises in simple realizations of the doublet-triplet splitting.
vev, and that direct mass terms $M_{ij} \mathbf{10}, \mathbf{10}_j$ are forbidden (which can easily be achieved by a global symmetry). In the opposite case, the light $L_i$ and $d^c_i$ are combinations of the corresponding components of $\mathbf{5}_i^{10}$ and $\mathbf{5}_i^{16}$, and many of the features described above are modified. In particular, the GUT-scale relation $M_d = M_e^T$ no longer hold $\mathbf{10}$, and neutrino masses receive both type I and type II seesaw contributions. Since we are interested in the more predictive type II case, we shall not consider this possibility any longer.

Let us close this section with some remarks about the seesaw scale. The observed frequency of atmospheric neutrino oscillations indicates that $M_\Delta$ in Eq. (5) should lie significantly below the GUT scale, $M_\Delta \lesssim 5 \times 10^{14}$ GeV. On the other hand, it is difficult to maintain gauge coupling unification with incomplete $SU(5)$ representations at an intermediate scale. Since $\mathbf{54} = \mathbf{24} + \mathbf{15} + \mathbf{15}$ under $SU(5)$, with $\Delta \subset \mathbf{15}$ ($\Delta \subset \mathbf{15}$), this suggests two possibilities (i) the whole $\mathbf{54}$ representation lies below the GUT scale ($M_\Delta = M_{54} \ll M_{GUT}$); (ii) the $\mathbf{24}^{54}$ component of $\mathbf{54}$ is split from the ($\mathbf{15}^{54}, \mathbf{15}^{54}$) pair ($M_\Delta = M_{15} \ll M_{24}$). The first possibility can be realized by generating an effective mass for the $\mathbf{54}$ from the non-renormalizable operator $\bar{\mathbf{16}} \mathbf{15} \mathbf{54} \mathbf{54}/\Lambda$, while forbidding a direct mass term by an appropriate symmetry. The second possibility requires a splitting between $M_{24}$ and $M_{15}$. This can be achieved e.g. by an additional coupling $\mathbf{54} \mathbf{45} \mathbf{45}$ involving a $\mathbf{45}$ Higgs representation with a GUT-scale vev along its $SU(5)$-singlet component, which will give a large mass to the $\mathbf{24}^{54}$ component of $\mathbf{54}$.

Notice that the requirement of perturbative gauge coupling unification constrains the values of $M_{15}$ and $M_{24}$. Indeed, while the unification of gauge couplings at $M_{GUT} \simeq 2 \times 10^{16}$ GeV is not spoiled by the addition of complete $SU(5)$ representations to the MSSM spectrum, their contribution to the beta function coefficients increases the value of the unified gauge coupling. Taking e.g. $M_{24} = 10M_{15} = 100M_1$ (and noting that Eq. (1) implies $M_i \propto m_{e_i}$, where the $M_i$ are the masses of the ($\mathbf{5}_i^{10}, \mathbf{5}_i^{16}$) pairs), the requirement of perturbative unification, $\alpha_{GUT} < 1$, sets a lower bound of $4 \times 10^{11}$ GeV on $M_{15}$. This bound can be lowered if the ($\mathbf{15}^{54}, \mathbf{15}^{54}$) and/or $\mathbf{24}^{54}$ multiplets are split, or if $M_{24} \gg 10M_{15}$.

### 3 The leptogenesis scenario

The class of $SO(10)$ models described in the previous section contains several heavy states that can contribute to leptogenesis through their decays. Let us first consider the components of the ($\mathbf{15}^{54}, \mathbf{15}^{54}$) multiplets, which as discussed above must be significantly lighter than $M_{GUT}$ (we shall address the case of the $\mathbf{24}^{54}$ multiplet in Section 5.2). Under $SU(3)_c \times SU(2)_L \times U(1)_Y$, the $\mathbf{15}$ representation decomposes as $\mathbf{15} = (\Sigma, Z, \Delta)$, where $\Sigma = (6,1)_{-2/3}$, $Z = (3,2)_{+1/6}$ and $\Delta = (1,3)_{+1}$. Since $\mathbf{10}, \mathbf{10}, \mathbf{54}, \mathbf{54} \supset \mathbf{5}_i^{10}, \mathbf{5}_j^{16} \mathbf{15}^{54}$, each of these states can decay into a pair of light matter fields. The supergraphs responsible for the asymmetries between these decays and the CP-conjugated decays are represented in Fig. 1 where, depending on the decaying state, $\mathbf{15}^{54}$, $\mathbf{15}^{54}$, $\mathbf{24}^{54}$, $\mathbf{5}_i^{10}$ and $\mathbf{5}_j^{10}$ should be replaced by their appropriate

---

3Note that no component of the $\mathbf{54}$ supermultiplet can mediate proton decay, so an intermediate-scale $\mathbf{54}$ or $\mathbf{15} + \mathbf{15}$ pair does not cause any phenomenological problem.
components. Note that there is no self-energy contribution to the CP asymmetry, since the interference of the self-energy diagram with the tree level diagram is proportional to the real combination of couplings $\text{Tr}(ff^*)^2$. If all three generations of $\overline{5}_i$’s were massless or degenerate in mass, the vertex contribution would be proportional to $\text{Im}[\text{Tr}(ff^*)] = 0$ and would therefore vanish as well.

However, this is not the case in our scenario, since the $(\overline{5}_i, \overline{5}_j)$ vector-like pairs are heavy with strongly hierarchical masses. Indeed, as follows from the second term in Eq. (1), their masses are determined by the same couplings as the charged lepton masses (up to possible corrections from non-renormalizable operators modifying the mass relation $M_d = M_T^T$): $M_i = h_i v^i_{16} \approx m_e (v^i_{16}/v^d_{16})$, where the $h_i$ are the eigenvalues of the matrix of couplings $h_{ij}$. Furthermore, the optical theorem tells us that only the $(\overline{5}_i, \overline{5}_j)$ pairs with $M_{15} > M_i + M_j$ can contribute to the CP asymmetry. The vertex contribution is therefore proportional to $\sum_{ij} c_{ij} \theta(M_{15} - M_i - M_j) \text{Im}[f_{ij}(f^* f^*)_{ij}]$, where the coefficients $c_{ij}$ account for the $M_i$-dependent loop function, and $\theta(x)$ is the Heaviside function. Thus, thanks to the presence of heavy states with hierarchical masses in the loop, the CP asymmetry does not vanish in our scenario. This is very different from the usual triplet leptogenesis scenarios [11], in which two different sets of couplings are needed in order to obtain a non-vanishing CP asymmetry.

Notice that (i) the three relevant $SU(5)$ couplings $\overline{5}_i \overline{5}_j \overline{5}_k \overline{5}_l$, $\overline{5}_i \overline{5}_j \overline{5}_k \overline{5}_l$, $\overline{5}_i \overline{5}_j \overline{5}_k \overline{5}_l$, and $\overline{5}_i \overline{5}_j \overline{5}_k \overline{5}_l$ are determined by one and the same matrix $f$; (ii) the masses of the $\overline{5}_i$’s are approximatively known, up to an overall scale, in terms of charged lepton masses. These properties, which follow from the underlying $SO(10)$ structure, make this model more predictive than usual leptogenesis scenarios.

3.1 The CP asymmetry in triplet decays

In this section, we compute the CP asymmetry in triplet decays. The relevant superpotential, inherited from the $SO(10)$ superpotential (1), reads:

$$W_{\Delta \bar{\Delta}} = M_{\Delta \bar{\Delta}} \text{Tr}(\Delta \bar{\Delta}) + \frac{1}{2} f_{ij} (L_i^T i \sigma_2 \Delta L_j - L_i^T i \sigma_2 \bar{\Delta} \bar{L}_j) + \frac{1}{2} (\sigma_u H_u^T i \sigma_2 \Delta H_u - \sigma_d H_d^T i \sigma_2 \Delta H_d),$$

(6)
where $H_u$ and $H_d$ are the light Higgs doublets, $\sigma_{u,d} \equiv \sigma \alpha_{u,d}^2$ ($\alpha_{u,d}$ being defined by $H_{d,u}^{10} = \alpha_{d,u} H_{d,u} + \ldots$) and

$$\Delta \equiv \frac{\sigma}{\sqrt{2}} \cdot \Delta = \begin{pmatrix} \Delta^+ / \sqrt{2} & \Delta''+ / \sqrt{2} \\ \Delta^0 & -\Delta^+ / \sqrt{2} \end{pmatrix}, \quad L_i = \begin{pmatrix} \nu_i \\ e_i \end{pmatrix}, \quad (7)$$

$$\bar{\Delta} \equiv \frac{\sigma}{\sqrt{2}} \cdot \bar{\Delta} = \begin{pmatrix} \bar{\Delta}''- / \sqrt{2} & \bar{\Delta}^0 \\ \bar{\Delta}^- / \sqrt{2} & -\bar{\Delta}^0 / \sqrt{2} \end{pmatrix}, \quad L_i = \begin{pmatrix} E_i^c \\ -N_i^c \end{pmatrix}. \quad (8)$$

The following superpotential terms will also be needed for the computation of the CP asymmetry:

$$W_{S,T} = f_{ij} L_i^c T \left( \frac{1}{2} \sqrt{\frac{3}{5}} S + \sqrt{\frac{1}{2}} T \right) i\sigma_2 L_j + \frac{1}{2} M_S S^2 + \frac{1}{2} M_T \text{Tr}(T^2), \quad (9)$$

where $S$ and $T$ are the $(1,1,0)$ and $(1,3,0)$ components of $24^{54}$, and $T \equiv \bar{\sigma} \cdot \bar{T} / \sqrt{2}$.

The chiral superfields $\Delta$ and $\bar{\Delta}$ describe two complex scalar fields, $\Delta_s$ and $\bar{\Delta}_s$, and a Dirac spinor field $\Psi_\Delta$. Since we are dealing with a supersymmetric theory, we only need to compute the CP asymmetry in the decays of one of the component fields, e.g. the scalar triplet $\Delta_s$. This field has four decay modes (neglecting phase space suppressed 3-body decays such as $\Delta_s \rightarrow \bar{L}c\bar{c} H_d$): it can decay into light leptons ($\Delta_s \rightarrow LL$), heavy sleptons ($\Delta_s \rightarrow \bar{L}c \bar{L}$), down-type Higgsinos ($\Delta_s \rightarrow \bar{H}_d \bar{H}_d$), and up-type Higgs bosons ($\Delta_s \rightarrow H_u H_u$). The corresponding decay widths are given by

$$\Gamma(\Delta_s \rightarrow aa) = \frac{M_\Delta}{32\pi} \lambda_a^2, \quad a = T, \bar{T}, \bar{H}_d, H_u, \quad (10)$$

with

$$\lambda_L^2 \equiv \sum_{i,j=1}^3 |f_{ij}|^2, \quad \lambda_{Lc}^2 \equiv \sum_{i,j=1}^3 K \left( \frac{M_i^2}{M_\Delta^2}, \frac{M_j^2}{M_\Delta^2} \right) |f_{ij}|^2, \quad \lambda_{H_{u,d}}^2 \equiv |\sigma_{u,d}|^2. \quad (11)$$

The kinematic factor $K$ for decays into $\bar{L}c$'s is

$$K(x_i, x_j) = \theta(1 - \sqrt{x_i} - \sqrt{x_j}) \sqrt{\lambda(1, x_i, x_j)}, \quad (12)$$

where $\lambda(x, y, z) \equiv x^2 + y^2 + z^2 - 2(xy + yz + xz)$.

The CP asymmetry in the decays of $\Delta_s$, $\Delta_s^*$ into light leptons is defined by

$$\epsilon_{\Delta_s \rightarrow LL} \equiv \frac{\Gamma(\Delta_s \rightarrow LL) - \Gamma(\Delta_s^* \rightarrow LL)}{\Gamma_{\Delta_s} + \Gamma_{\Delta_s^*}}, \quad (13)$$

where $\Gamma_{\Delta_s} = M_\Delta \sum_a \lambda_a^2 / (32\pi)$ is the total $\Delta_s$ decay rate. Unitarity and CPT imply $\Gamma_{\Delta_s} = \Gamma_{\Delta_s^*}$. One can easily check that the asymmetries in the decays $\Delta_s \rightarrow \bar{H}_d \bar{H}_d$ and $\Delta_s \rightarrow H_u H_u$ vanish at the one loop order, as they involve a single coupling $\sigma_d$ or $\sigma_u$. As a result, the equality $\Gamma_{\Delta_s} = \Gamma_{\Delta_s^*}$ reduces to $\Gamma(\Delta_s^* \rightarrow LL) + \Gamma(\Delta_s^* \rightarrow \bar{L}c\bar{L}c) = \Gamma(\Delta_s \rightarrow LL) + \Gamma(\Delta_s \rightarrow \bar{L}c\bar{L}c)$. 

6
\[ \Gamma(\Delta_s \rightarrow L \bar{L}) + \Gamma(\Delta_s \rightarrow \tilde{L} \tilde{L}^c), \] i.e. the CP asymmetries in \( \Delta_s \) decays into light leptons and into heavy sleptons are exactly opposite. This allows one to define

\[ \epsilon_{\Delta} \equiv \frac{2}{3} \epsilon_{\Delta_s \rightarrow \tilde{L} \tilde{L}^c} = -2 \epsilon_{\Delta_s \rightarrow L \bar{L}}, \tag{14} \]

where the factor of 2 accounts for the fact that two antileptons are produced in each \( \Delta_s \) decay. Furthermore, supersymmetry ensures that the CP asymmetries of all components of the \( \Delta, \Delta \) supermultiplets are the same:

\[ \epsilon_{\Delta s} = \epsilon_{\bar{\Delta} s} = \epsilon_{\Psi} \equiv \epsilon_{\Delta}. \tag{15} \]

The Feynman diagrams relevant to the computation of \( \epsilon_{\Delta} \) are shown in Fig. 2. For arbitrary masses \( M_{\Delta}, M_S, M_T \) and \( M_i \) (\( i = 1, 2, 3 \)), one obtains:

\[ \epsilon_{\Delta} = \frac{1}{16 \pi} \sum_{R=S,T} \sum_{k,l=1}^3 c_R F \left( \frac{M_R^2}{M_{\Delta}}, \frac{M_k^2}{M_{\Delta}} \right) \frac{\text{Im}[f_{kl}(f^* f^* )_{kl}]}{32 \pi \Gamma_{\Delta_s}/M_{\Delta}}, \tag{16} \]

where \( c_S = 3/5 \) and \( c_T = 1 \) are \( SU(5) \) Clebsch-Gordan coefficients, and the loop function \( F \) reads:

\[ F(x, x_k, x_l) = \theta(1 - \sqrt{x_k} - \sqrt{x_l}) \sqrt{x} \ln \left[ 1 + 2x - x_k - x_l + \sqrt{\lambda(1, x_k, x_l)} \right] \frac{1 + 2x - x_k - x_l}{1 + 2x - x_k - x_l - \sqrt{\lambda(1, x_k, x_l)}}. \tag{17} \]

It is instructive to consider some particular cases. In the case \( 2M_3 < M_{\Delta}, \) all terms in the sum over \( k,l \) in Eq. (16) contribute to the asymmetry, and they add up to zero in the limit of massless \( L_i^c \)'s (\( M_i/M_{\Delta} \to 0 \)). This is simply due to the fact that \( \epsilon_{\Delta} \propto \text{Im} [\text{Tr}(f f^* f f^*)] \) in this limit, as discussed previously. In the case \( 2M_1 < M_{\Delta} < M_1 + M_2, \) only the term \( k = l = 1 \) contributes and \( F \) is maximal for \( M_1 = 0, \) while it is reduced by about a factor of 2 (4) for \( 2M_1/M_{\Delta} = 0.87 \) (0.97) and \( M_{S,T}/M_{\Delta} \geq 3.\) If in addition \( M_S = M_T \equiv M_{24}, \) Eq. (16) simplifies to, for \( M_1 \ll M_{\Delta}: \)

\[ \epsilon_{\Delta} \approx \frac{1}{10 \pi} \frac{M_{24}^2}{M_{\Delta}} \ln \left( 1 + \frac{M_{24}^2}{M_{24}^2} \right) \frac{\text{Im}[f_{11}(f^* f^* )_{11}]}{\lambda_L + \lambda_{L^c}^2 + \lambda_{H_u}^2 + \lambda_{H_d}^2}, \tag{18} \]

where \( \lambda_{L_i^c} \equiv |f_{11}|. \)
3.2 Other CP asymmetries

The CP asymmetries in the decays of the other components of the $(15, \overline{15})$ multiplets can be computed along the same lines, taking into account the following differences with the triplet case (we consider the scalar components of the $\Sigma$ and $Z$ multiplets for definiteness):

- $\Sigma_s$ and $Z_s$ cannot decay into Higgs fields, since the colour triplets contained in the 10 Higgs multiplets (to which the $\Sigma$ and $Z$ fields couple via the $SO(10)$ operator $10 \otimes 10 \otimes 54$) must have GUT-scale masses in order to suppress the dangerous $D = 5$ proton decay operators. Thus, they have only two 2-body decay modes, $\Sigma_s \to \overline{d}d$ and $\Sigma_s \to \overline{D}D$ (respectively $Z_s \to \overline{L}d$ and $Z_s \to \overline{L}\overline{D}$), with opposite CP asymmetries:

$$\epsilon_{\Sigma} \equiv 2\epsilon_{\Sigma_s \to \overline{d}d} = -2\epsilon_{\Sigma_s \to \overline{D}D},$$

$$\epsilon_{Z} \equiv \epsilon_{Z_s \to L\overline{d}} = -\epsilon_{Z_s \to L\overline{D}}.$$  

- $\epsilon_{\Sigma}$ and $\epsilon_{Z}$ are given by a similar expression to Eq. (16), with $M_\Delta$ replaced by $M_{\Sigma}$ and $M_Z$, respectively; however the $24^{54}$ components in the loop and the Clebsch-Gordan coefficients are not the same. Furthermore, the total decay rates $\Gamma_{\Sigma_s}$ and $\Gamma_{Z_s}$ differ from $\Gamma_{\Delta_s}$ due to the absence of Higgs decay modes. For $|\sigma| \gg |f_{ij}|$ this results in a significant enhancement of $\epsilon_{\Sigma}$ and $\epsilon_{Z}$ with respect to $\epsilon_{\Delta}$.

It is important to notice that, among the components of the $(15, \overline{15})$ multiplets, only the $SU(2)_L$ triplets $(\Delta, \overline{\Delta})$ have $(B - L)$-violating interactions (in the limit in which $M_{GUT}$-suppressed interactions are neglected)$^4$. Hence the decays of the $(\Sigma, \overline{\Sigma})$ and $(Z, \overline{Z})$ fields cannot generate a sizeable $B - L$ asymmetry by themselves. Still they produce asymmetries in the number densities of the species $L_i$, $d_i$, $L_i^c$ and $D_i$, which can affect the dynamical evolution of the $B - L$ asymmetry generated in the decays of the $(\Delta, \overline{\Delta})$ components.

As is similar statement can be made about the decays of the components of the $24^{54}$. Feynman diagrams analogous to the ones shown in Fig. 2$^2$ generate a CP asymmetry between, e.g., the decay $S_s \to \overline{L}L^c$ and the CP-conjugated decay. However, since all renormalizable interactions of the $24^{54}$ components respect $B - L$, their decays cannot generate a $B - L$ asymmetry by themselves. Still they affect leptogenesis by producing asymmetries in the number densities of the species $L_i$, $L_i^c$, $d_i^c$ and $D_i$. In the following, we assume $M_{24} \gg M_\Delta$, so that either the components of the $24^{54}$ are not efficiently produced in the thermal bath after reheating, or the asymmetries generated in their decays have been washed out before the components of the $(\Delta, \overline{\Delta})$ supermultiplets decay.

Finally, the right-handed neutrinos are a source of $B - L$ violation. In our scenario, they do not have standard Dirac couplings, but they couple to the heavy leptons through the superpotential terms $y_u\nu^cL^{16}H_u^{10}$ and $y_d\nu^cL^cH_d^{16}$. If kinematically allowed,
the associated two-body decays will generate a $B - L$ asymmetry stored in the heavy
(s)leptons, before these decay to light species. In the following, we assume that the
right-handed neutrinos are heavier than the $(\Delta, \bar{\Delta})$ fields, so that we can neglect their
contribution to the final baryon asymmetry – either because they are not efficiently
produced after reheating, or because the asymmetries generated in their decays have
been washed out before the triplets decay. This assumption is analogous to the one
usually done in type I leptogenesis, where the contribution of the next-to-lightest right-
handed neutrino to the lepton asymmetry is neglected. Also, the decays of the heavy
lepton fields into right-handed neutrinos, which would introduce an additional source
of $B - L$ violation in our scenario, are kinematically forbidden.

3.3 Dependence of the CP asymmetry on the light neutrino mass parameters

The CP asymmetry in triplet decays, Eq. (16), depends on the heavy lepton masses
and couplings, which in turn are related to the low-energy lepton parameters. Indeed,
the superpotential terms $h_{ij}^{16}\bar{10}_i \bar{10}_j^{16}$ yield the following GUT-scale mass relations:

$$M_i = m_{e_i} \frac{v_1^{16}}{v_d^{16}} = m_{d_i} \frac{v_1^{16}}{v_d^{16}},$$

(21)

where $v_1^{16}$ is the vev of the $SU(5)$-singlet component of $16$. Furthermore, it is possible
to choose a $\bar{10}_i$ basis in which the charged lepton and heavy matter mass matrices are
simultaneously diagonal, and in this basis:

$$f = \frac{2M_\Delta}{\lambda_{H_u}^2 v^2 \sin^2 \beta} U^\dagger \text{Diag} (m_1, m_2, m_3) U^\dagger,$$

(22)

where the $m_i$ are the light neutrino masses and $U$ is the PMNS mixing matrix. The
non-renormalizable operators needed to correct the GUT-scale relations $m_\mu = m_s$ and
$m_e = m_d$ will in general modify Eq. (21), but one can neglect this effect for an order-of-
magnitude estimate of the heavy lepton masses. With the $m_{e_i}$ evaluated at the GUT
scale, and assuming that the light Higgs doublet $H_d$ contains a large $H_{16}^d$ component
(i.e. $v_d^{16} \sim v_d$), one obtains:

$$(M_1, M_2, M_3) \sim (2 \times 10^{11}, 4 \times 10^{13}, 7 \times 10^{14}) \text{ GeV} \left(\frac{\tan \beta}{10}\right) \left(\frac{v_1^{16}}{10^{16} \text{ GeV}}\right).$$

(23)

Gauge coupling unification favours values of $v_1^{16}$ close to $M_{GUT}$, unless the rank of
$SO(10)$ is broken at the GUT scale by the vev of a different Higgs multiplet (e.g. an
extra $16$), in which case $v_1^{16}$, hence the $M_i$, can be significantly smaller.

In the following, we assume that the situation $M_1 \ll M_\Delta < M_1 + M_2$ is realized.\footnote{If instead $M_1 + M_2 < M_\Delta$, $\epsilon_\Delta$ receives extra contributions, opening additional possibilities to achieve successful leptogenesis that could be analyzed along the same lines.}

The CP asymmetry $\epsilon_\Delta$ is then given by Eq. (18), which can be rewritten as:

$$\epsilon_\Delta \simeq \frac{1}{10\pi} g \left(\frac{M_2^2}{M_\Delta^2}\right) \lambda_{L}^{2} \lambda_{L}^{2} \lambda_{L}^{2} \lambda_{L}^{2} \lambda_{H_u}^{2} \lambda_{H_u}^{2} \lambda_{H_d}^{2} \lambda_{H_d}^{2} \frac{\text{Im}[m_{11}(m^{*}mm^{*})_{11}]}{m_{11}^3},$$

(24)
where $g(x) \equiv \sqrt{x} \ln [1 + 1/x]$, $m \equiv U^* \text{Diag} (m_1, m_2, m_3) U^\dagger$ and $\overline{m}^2 \equiv \sum_i m_i^2$. In writing Eq. (21), we assumed the absence of a mismatch between the mass eigenstate bases of charged leptons and of the heavy lepton fields, as in Eq. (22). Such a mismatch may arise from the corrections needed to account for $m_{d_i} \neq m_{e_i}$, but Eq. (21) is still a good approximation if the unitary matrix that describes this mismatch is characterized by small mixing angles (compared with the uncertainties on the PMNS angles).

The factor in Eq. (21) that explicitly depends on the light neutrino mass parameters reads:

$$\frac{\text{Im}[m_{11}(m^* m m^*)_{11}]}{m^1} = -\frac{1}{m^1} \left\{ c_{13}^2 c_{12}^2 s_{12}^2 \sin(2\rho) m_1 m_2 \Delta m_{21}^2 + c_{13}^2 s_{13} c_{12}^2 2(\rho - \sigma) m_1 m_2 \Delta m_{31}^2 - c_{13}^2 s_{13}^2 c_{12}^2 2(2\sigma) m_2 m_3 \Delta m_{32}^2 \right\}, \quad (25)$$

where $c_{ij} \equiv \cos \theta_{ij}$, $s_{ij} \equiv \sin \theta_{ij}$, $\Delta m_{ij}^2 \equiv m_j^2 - m_i^2$, and we adopted the parametrization $U_{ei} = (c_{13} c_{12}^{e} i \rho, c_{13} s_{12}^{e}, s_{13}^{e} e^{i\sigma})$, in which $\rho$ and $\sigma$ are the two Majorana-type CP-violating phases to which neutrinoless double beta decay is sensitive. Indeed, neutrinoless double beta decay depends on the effective Majorana mass $m_{\nu e c} = |\sum_i m_i U_{ei}^2|$. Some comments about Eq. (21) are in order. First, to the extent discussed above, the CP asymmetry is predicted in terms of the light neutrino parameters once $\tan \beta$ and the flavour-blind parameters $\lambda_L$, $\lambda_{H_u}$, $\lambda_{H_d}$ and $M_\Delta/M_24$ are specified ($\lambda_{L_\nu}$ is not an independent parameter, since it is related to $\lambda_L$ by $\lambda_{L_\nu} = \lambda_L |m_{11}|/\overline{m}$). This is a noticeable difference with leptogenesis in the standard type I and type II seesaw mechanisms. Second, the CP asymmetry depends on the same two CP-violating phases as neutrinoless double beta decay; observing CP violation in neutrino oscillations would not be enough to test the validity of the present scenario. Finally, it is very sensitive to the yet unknown value of $\theta_{13}$ and to the type of mass spectrum.

To illustrate this point, let us look for the parameter values that maximize the quantity in Eq. (25). If $\theta_{13}$ is close to its present experimental upper bound, $\sin^2 \theta_{13} \approx 0.05$, the maximal value of Eq. (25) is obtained for a normal neutrino mass hierarchy with $m_1 \approx m_2 \approx 0.01$ eV and $\rho = 0$, $\sigma = \pi/4$, and it is given by $\approx \frac{\sin^2 \theta_{13}}{c_{13}^2 s_{13}^2 \sqrt{\Delta m_{21}^2 / \Delta m_{32}^2}}$. A more precise estimate gives $9.2 \times 10^{-3}$ for $m_1 = 0.016$ eV and $s_{13}^2 \approx 0.05$. For $\sin^2 \theta_{13} \lesssim 0.01$, the maximum value of Eq. (25) is obtained for an inverted hierarchy with $m_3 = 0$ and $\rho = -\pi/4$, and it is given by $\approx \frac{\sin^2 \theta_{13}}{c_{12}^2 \Delta m_{21}^2 / (4|\Delta m_{32}^2|)}$. A more precise estimate gives $1.7 \times 10^{-3}$ for $\theta_{13} = 0$. Taking $\lambda_{L_\nu}^2 \gg \lambda_{H_u}^2, \lambda_{H_d}^2$ and choosing $M_{24} \approx M_\Delta/2$ in order to maximize the loop function $g(M_{24}^2/M_\Delta^2)$, one obtains, for the above two sets of light neutrino mass parameters:

$$\epsilon_\Delta \approx 2.2 \times 10^{-4} \lambda_L^2 \quad \text{(maximum $\theta_{13}$)}, \quad (26)$$

$$\approx 3.4 \times 10^{-5} \lambda_L^2 \quad \text{(vanishing $\theta_{13}$)}, \quad (27)$$

where the value of $\lambda_L^2$ is bounded by perturbativity (requiring $|f_{ij}| \leq 1$, one can take $\lambda_L^2$ as large as 5 for a normal hierarchical spectrum). In the above estimates, we used $|\Delta m_{31}^2| = 2.4 \times 10^{-3}$ eV$^2$, $\Delta m_{21}^2 = 7.6 \times 10^{-5}$ eV$^2$ and $\sin^2 \theta_{12} = 0.32$ [12]. The asymmetry, maximized with respect to the CP-violating phases and to $M_\Delta/M_{24}$, is plotted in Fig. 3 as a function of the lightest neutrino mass and $\theta_{13}$, for the cases of normal and inverted hierarchy.
Figure 3: The CP-asymmetry $\epsilon_\Delta$ in units of $\lambda_L^2$ as function of the lightest neutrino mass and of $\sin^2 \theta_{13}$ in the cases of normal and inverted hierarchy of light neutrino masses. The parameters $|\Delta m_{32}^2|$, $\Delta m_{21}^2$ and $\theta_{12}$ are chosen as specified in the text. The asymmetry is maximized with respect to the CP-violating phases and to the ratio $M_\Delta/M_{24}$.

Eq. (26), together with Eq. (29), shows that an efficiency factor as small as $10^{-5} - 10^{-4}$ is sufficient for successful baryogenesis if $\epsilon_\Delta$ is close to its maximum value. A quantitative estimate of $\eta$ requires in general the numerical resolution of the Boltzmann equations, and is beyond the scope of the present paper. However, the region of the parameter space which leads to a large efficiency factor can be discussed analytically, as we do in Section 4.

4 Dynamical evolution of the $B-L$ asymmetry

As discussed in Section 3.2, in the limit in which one neglects $M_{GUT}$-suppressed interactions, the only source for the $B-L$ asymmetry in our scenario are the out-of-equilibrium decays of the components of the $(\Delta, \bar{\Delta})$ supermultiplets. In principle, the $M_{GUT}$-suppressed interactions could be relevant for the subsequent decays of $L_1^c$ and $D_1$, which would then violate $B-L$ and affect the final baryon asymmetry. However, it turns out that the dominant decay modes of $L_1^c$ and $D_1$ preserve $B-L$ as long as the light Higgs doublet $H_d$ contains a non-negligible $H_{d_{10}}$ component, and we shall assume that this is the case in the following.

---

6A numerical analysis of the efficiency has been performed in Refs. [13, 14] for the simpler case of the Standard Model augmented with an $SU(2)_L$ Higgs triplet, and in Ref. [15] for its supersymmetric extension.

7Unless $\alpha_d$ is very small, the (B-L)-conserving decay modes, even though suppressed by small Yukawa couplings, dominate over the (B-L)-violating ones mediated by the heavy right-handed neutrinos or by the
As the $B-L$ asymmetry originates from the triplet decays, we expect that a good estimate of the efficiency can be obtained by considering the decays and interactions of the $(\Delta, \overline{\Delta})$ fields only, as we do below in order to simplify the discussion. We shall check in Section 4.3 that including all components of the $(15, \overline{15})$ multiplets in the analysis does not affect our conclusions. We therefore define the efficiency factor $\eta$ as follows:

$$n_{B-L} = \eta \epsilon_{\Delta} \left[ \frac{n_{eq}^{\Delta} + \Delta_{\ell}^{s}}{s} + \frac{n_{eq}^{\overline{\Delta}} + \Delta_{\ell}^{s}}{s} + \frac{n_{eq}^{\Psi_{\Delta}} + \overline{\Psi}_{\Delta}}{s} \right]_{T \gg M_{\Delta}}$$

where $s = (2\pi^{2}/45)g_{*}T^{3}$ is the entropy density. After conversion of the $B-L$ asymmetry by the $(B + L)$-violating sphalerons, the final baryon asymmetry reads:

$$\eta \epsilon_{\Delta} = 7.62 \times 10^{-3}$$

where we have used $g_{S} = g_{*S}(MSSM) + g_{*S}(5_{1}, \overline{5}_{1}) = 266.25$ (we assume as in the previous section that a single $(5_{10}, \overline{5}_{10})$ pair is lighter than the triplet, namely $M_{1} < M_{1} + M_{2}$). The observed value of the baryon-to-entropy ratio, $(n_{B}/s)_{WMAp} = (8.82 \pm 0.23) \times 10^{-11}$ [16], requires $\eta \epsilon_{\Delta} \simeq 10^{-8}$.

4.1 Boltzmann equations

Since the interaction rates and CP asymmetries involving different components of the $(\Delta, \overline{\Delta})$ superfields are related by supersymmetry, one can obtain a valid estimate of $\eta$ by considering solely the scalar triplet $\Delta_{s}$. The problem of computing $\eta$ is then similar to the case of the Higgs triplet extension of the Standard Model studied in Ref. [13], with however two important differences: our scenario involves another species carrying lepton number to which $\Delta_{s}$ can decay ($\tilde{L}_{i}^{c}$), and there is no CP asymmetry in the Higgs/Higgsino decay channels. The relevant Boltzmann equations (written for simplicity in the $a_{d} = 0$ case, in which $\Delta_{s}$ has no Higgsino decay mode) read:

$$sH \frac{d\Sigma_{\Delta_{s}}}{dz} = -\gamma_{D} \left( \frac{\Sigma_{\Delta_{s}}}{\Sigma_{eq}} - 1 \right) - 2\gamma_{A} \left( \left( \frac{\Sigma_{\Delta_{s}}}{\Sigma_{eq}} \right)^{2} - 1 \right),$$

$$sH \frac{d\Delta_{s}}{dz} = -\gamma_{D} \left( \frac{\Delta_{s}}{\Sigma_{\Delta_{s}}} + B_{L} \frac{\Delta_{L}}{Y_{L}^{eq}} - B_{H_{u}} \frac{\Delta_{H_{u}}}{Y_{H_{u}}^{eq}} - B_{L_{i}} \frac{\Delta_{L_{i}}}{Y_{L_{i}}^{eq}} \right),$$

$$sH \frac{d\Delta_{L}}{dz} = \gamma_{D} \epsilon_{\Delta} \left( \frac{\Sigma_{\Delta_{s}}}{\Sigma_{eq}} - 1 \right) - 2\gamma_{D} B_{L} \left( \frac{\Delta_{L}}{Y_{L}^{eq}} + \frac{\Delta_{\Delta_{s}}}{\Sigma_{eq}} \right),$$

$$sH \frac{d\Delta_{H_{u}}}{dz} = -2\gamma_{D} B_{H_{u}} \left( \frac{\Delta_{H_{u}}}{Y_{H_{u}}^{eq}} - \frac{\Delta_{\Delta_{s}}}{\Sigma_{eq}} \right),$$

$$sH \frac{d\Delta_{L_{i}}}{dz} = \gamma_{D} \epsilon_{\Delta} \left( \frac{\Sigma_{\Delta_{s}}}{\Sigma_{eq}} - 1 \right) - 2\gamma_{D} B_{L_{i}} \left( \frac{\Delta_{L_{i}}}{Y_{L_{i}}^{eq}} - \frac{\Delta_{\Delta_{s}}}{\Sigma_{eq}} \right).$$

colour triplets. For instance, the dominant decay mode of $\tilde{L}_{i}$ (respectively $\tilde{D}_{1}$) is $\tilde{L}_{i} \rightarrow \tilde{e} H_{d}$ (respectively $\tilde{D}_{1} \rightarrow Q H_{d}$) over most of the parameter space. It is interesting to note that natural realizations of the doublet-triplet splitting seem to require $a_{d} \neq 0$ in our scenario.
where \( z = M_\Delta/T, Y_X = n_X/s, \Delta_X \equiv Y_X - Y_{\overline{X}} \) is the asymmetry stored in the species \( X \), and \( \Sigma_{\Delta_x} \equiv Y_{\Delta_x} + Y_{\overline{\Delta_x}} \) is the total density of \( \Delta_s \) and \( \Delta_s^* \); \( \gamma_D = s \Sigma_{\Delta_x} \Gamma_{\Delta_x}/K_1(z)/K_2(z) \) (where \( K_1(z) \) and \( K_2(z) \) are modified Bessel functions) and \( \gamma_A \) are the space-time densities of the \( \Delta_s, \Delta_s^* \) decays and of the gauge scatterings \( \Delta_s \Delta_s^* \rightarrow \text{lighter particles} \), respectively; \( B_L, B_{H_u} \) and \( B_{L_L} \) are the branching ratios of \( \Delta_s \) decays into \( \overline{T}T, H_uH_u \) and \( L_L \), respectively. Note that we included only decays, inverse decays and gauge scatterings in Eqs. (30) to (34). In particular, we omitted the triplet-mediated \( \Delta_L = 2 \) scatterings \( LL \leftrightarrow H_uH_u^* \) and \( LH_u \leftrightarrow \overline{T}H_u^* \), which due to the smallness of neutrino masses are much slower that the expansion of the Universe (except for very large values of \( M_\Delta \)). Scatterings such as \( L_L \leftrightarrow H_uH_u \) (\( \Delta_L = 2 \), triplet-mediated) and \( L_L \leftrightarrow \overline{T}L \) (\( \Delta_L = 0 \), triplet- or \( S, T \)-mediated) are even slower in the region of the parameter space that we shall consider below. We also neglected the washout due to the inverse decays of \( S \) and \( T \), which are Boltzmann suppressed at \( T \sim M_\Delta \) since \( M_\Delta \ll M_{24} \).

An important point to notice is that the Boltzmann equations (31) to (34) are not linearly independent. As a result, the following combination of asymmetries is preserved:

\[
2\Delta_{\Delta_s} - \Delta_L + \Delta_{H_u} + \Delta_{L_L} = 0. \tag{35}
\]

This relation generalizes the sum rule of Ref. [13].

### 4.2 Conditions for an order one efficiency

In triplet leptogenesis, an order one efficiency can be obtained even though the total triplet decay rate is larger than the expansion rate of the universe, provided that one of the decay channels is out of equilibrium and sufficiently decoupled from the fast channel(s). This has been first pointed out in Ref. [13], where the case of the Standard Model augmented with a scalar triplet has been studied. A similar conclusion can be reached in our scenario, with the role of the slow decay mode played by \( \Delta_s \rightarrow L_L \overline{L}_L \), as we now discuss.

Let us first consider the out-of-equilibrium conditions for the various decay channels. The condition for the decay \( \Delta_s \rightarrow aa \ (a = \overline{T}, H_u, \overline{L}_L) \) to be out of equilibrium at \( T = M_\Delta \) is \( K_a \equiv \Gamma(\Delta_s \rightarrow aa)/H(M_\Delta) \ll 1 \). Since \( \Gamma(\Delta_s \rightarrow aa) = \lambda_a^2 M_\Delta/(32\pi) \) and \( H(T) = 1.66\sqrt{g_\ast} T^2/M_P \) (with \( g_\ast = g_{sS} = 266.25 \) during leptogenesis), this condition translates into:

\[
\lambda_a \ll 10^{-2} \sqrt{\frac{M_\Delta}{10^{12} \text{GeV}}}, \tag{36}
\]

where we recall that \( \lambda_L = \sqrt{\text{Tr}(f f^*)}, \lambda_{H_u} = |\sigma_u|, \) and \( \lambda_{L_L} = |f_{11}|. \) Due to \( \tilde{m}^2 \equiv \sum_i m_i^2 \propto \lambda_L^2 \lambda_{H_u}/M_\Delta^2 \), the product \( K_L K_{H_u} \) is controlled by the scale of neutrino masses:

\[
K_L K_{H_u} \simeq \frac{220}{\sin^4 \beta} \left( \frac{\tilde{m}}{0.05 \text{eV}} \right)^2. \tag{37}
\]

Hence, at least one of the two channels \( \Delta_s \rightarrow \overline{T}T \) and \( \Delta_s \rightarrow H_uH_u \) (easily both) must be in equilibrium. Moreover, we have \( \lambda_{L_L} < \lambda_L \) (with \( \lambda_{L_L} \lesssim 0.2 \lambda_L \) for hierarchical light neutrino masses). The candidate out-of-equilibrium decay is therefore \( \Delta_s \rightarrow L_L \overline{L}_L. \) As
we are going to see, a large efficiency can be reached in the region \( \lambda_{L_i} \ll \lambda_{H_u}, \lambda_L \), in which triplet decays into leptons and Higgs are in equilibrium\(^8\) (\( K_{H_u}, K_L \gtrsim 1 \)), while decays into \( \tilde{L}_i \tilde{L}_1 \) are not (\( K_{\tilde{L}_i} \ll 1 \)).

Gauge scatterings \((\Delta_s \Delta_s^* \rightarrow \text{lighter particles})\) first create an equilibrium population of triplets and antitriplets. During leptogenesis, this population is kept close to thermal equilibrium by decays and inverse decays, thanks to Eq. (37) (the fact that \( \gamma_A < \gamma_D \) for \( T < M_\Delta \) allows the triplets to decay before annihilating\(^{[13]}\)). In spite of this, a large asymmetry \( \Delta_{L_i} \) can develop without being washed out, due to the fact that \( K_{\tilde{L}_i} \ll 1 \). This implies an asymmetry between the abundances of triplets and antitriplets, which is then transferred to \( \Delta_L \) and \( \Delta_{H_u} \) through their decays. After all triplets have decayed, and before the \( \tilde{L}_i \)'s decay, we end up with:

\[
Y_{B-L} = \Delta_{\tilde{L}_i} - \Delta_L = -\Delta_{H_u},
\]

where we made use of the sum rule (35). In order to create a large \( B-L \) asymmetry, a large \( \Delta_{H_u} \) is thus needed. We can check that a large \( \Delta_{H_u} \) indeed forms, and relate its size to \( \Delta_{\tilde{L}_i} \) by noticing that the combinations

\[
\frac{\Delta_{L_i}}{Y_{eq}^{\Sigma_{\Delta_s}}} + \frac{\Delta_{u}}{\Sigma_{\Delta_s}^{eq}} \quad \text{and} \quad \frac{\Delta_{H_u}}{Y_{eq}^{\Sigma_{\Delta_s}}} - \frac{\Delta_{L_i}}{\Sigma_{\Delta_s}^{eq}},
\]

which multiply \((-2\gamma_DB_L)\) in Eq. (32) and \((-2\gamma_DB_{H_u})\) in Eq. (33), respectively, are forced to vanish due to \( K_{\tilde{L}_i}, K_{H_u} \gtrsim 1 \). One can check that this is still the case after most triplets have decayed. Then, using again the sum rule (35), we obtain:

\[
Y_{B-L} \simeq \frac{Y_{eq}^{\Sigma_{\Delta_s}}}{Y_{eq}^{\Sigma_{\Delta_s}} + Y_{eq}^{\Sigma_{\Delta_s}}} \Delta_{\tilde{L}_i} = \frac{4}{7} \Delta_{\tilde{L}_i}.
\]

To estimate \( \Delta_{\tilde{L}_i} \), let us note that, in the limit \( \gamma_A = B_{L_i} = 0 \) (remember that \( \gamma_A \ll \gamma_D \) during leptogenesis), Eqs. (30) and (34) give \( \Delta_{\tilde{L}_i} = \epsilon \Delta \Sigma_{\Delta_s}^{eq} (T \gg M_\Delta) \). Taking into account the effect of gauge scatterings and the washout by inverse decays introduces an order one factor \( \eta_0 < 1 \) such that

\[
\Delta_{\tilde{L}_i} = \eta_0 \epsilon \Delta \Sigma_{\Delta_s}^{eq} (T \gg M_\Delta).
\]

The precise value of \( \eta_0 \) must be determined by solving numerically the complete set of Boltzmann equations, but the fact that \( \gamma_A \ll \gamma_D \) and \( K_{\tilde{L}_i} \ll 1 \) guarantees that it cannot be much smaller than 1. We conclude that an order one efficiency \( \eta \simeq 4\eta_0/7 \) can be reached in the region of parameter space where \( \lambda_{L_i} \ll \lambda_{H_u}, \lambda_L \). This region will be identified more precisely in Section 4.4. Although we set \( \alpha_d = 0 \) in the Boltzmann equations for simplicity, this value of \( \eta/\eta_0 \) is valid for non-vanishing values of \( \alpha_d \) such that \( K_{H_d} \ll 1 \) (let us recall that a non-negligible \( \alpha_d \) guarantees that the dominant \( \tilde{L}_i \) decay modes preserve \( B-L \)). For values of \( \alpha_d \) such that \( K_{H_d} \gtrsim 1 \), one obtains \( \eta \simeq 7\eta_0/10 \).

\(^8\)One may wonder what would happen if either \( K_{H_u} \ll 1 \) or \( K_L \ll 1 \). If \( K_{H_u} \ll 1 \), only a small \( \Delta_{H_u} \) is generated via Eq. (33), resulting in a suppressed \( Y_{B-L} \) according to Eq. (38). If \( K_L \ll 1 \), one has in particular \( \lambda_L \ll \lambda_{H_u} \) and this implies a strong suppression of \( \epsilon_{\Delta} \) according to Eq. (24).
Let us add a comment on lepton flavour effects. In the above, we worked in the "one-flavour approximation", i.e. we assumed a single generation of (light) leptons. Including the lepton flavours would greatly complicate the above discussion; however we can estimate their potential effect by assuming $\Delta_{\ell}\approx \Delta_{\mu}\approx \Delta_{\tau}$, which is not unreasonable given the mild hierarchy between the decay rates $\Delta \to T_iT_j$. In this case, the "flavoured" Boltzmann equations reduce to Eqs. (30) - (34) with $Y_L^\text{eq}$ replaced by $3Y_L^\text{eq}$. This leads to an efficiency $\eta \approx 4\eta_0/3$, to be compared with $\eta \approx 4\eta_0/7$ in the one-flavour case (keeping in mind that the value of $\eta_0$ itself depends on whether flavour effects are taken into account or not).

4.3 Effect of the other components of $(15, \overline{15})$

Let us now discuss how the above results are modified when all components of the $(15, \overline{15})$ multiplets are taken into account. In this case, Eqs. (30) - (34) are replaced with a larger number of Boltzmann equations describing the evolution of $\Sigma_{\Delta}, \Sigma_{\Sigma}, \Sigma_{Z_s}$ and of the asymmetries stored in the species $\Delta_{s}, \Sigma_{s}, Z_s, L, d^c, H_u, L_1^c$ and $D_1$. Instead of the sum rule (35), we now find two relations between asymmetries:

\begin{align}
2\Delta_{\Sigma_s} + \Delta_{Z_s} + \Delta_{D_{1i}} - \Delta_{d^c} &= 0, \\
2\Delta_{\Sigma_s} + \Delta_{Z_s} + \Delta_{L_{1i}} = -\Delta_{L} + \Delta_{H_u} &= 0.
\end{align}

The first relation tells us that $\Delta_{D_{1i}} = \Delta_{d^c}$ after all components of the $(15, \overline{15})$ have decayed, hence the coloured species do not contribute to the $B-L$ asymmetry. This is consistent with the fact that all relevant interactions preserve the baryon number. The second relation is analogous to Eq. (35). Furthermore, the combinations of asymmetries (39) are still driven to zero by the Boltzmann equations. As a result Eq. (40) still holds, but now both the $($asymmetries$)$ and the $($asymmetries$)$ decays contribute to $\Delta_{L_{1i}}$. In the limit where $\gamma_{\Delta}^A = \gamma_{\Sigma}^Z = B(\Delta \to \bar{L}_1\bar{L}_1) = B(Z \to \bar{L}_1\bar{D}_1) = 0$, one obtains $\Delta_{L_{1i}} = \epsilon_\Delta \Sigma_{\Delta}^\text{eq}_*(T \gg M_\Delta) + \epsilon_Z \Sigma_{\Sigma}^\text{eq}_*(T \gg M_Z) = (\epsilon_\Delta + 2\epsilon_Z) \Sigma_{\Delta}^\text{eq}_*(T \gg M_\Delta)$. Taking into account the effect of gauge scatterings and the washout by inverse decays introduces an order one factor $\eta_0 < 1$ such that

\begin{align}
\Delta_{L_{1i}} &= \eta_0(\epsilon_\Delta + 2\epsilon_Z) \Sigma_{\Delta}^\text{eq}_*(T \gg M_\Delta),
\end{align}

where, as discussed in Section 3.2, $\epsilon_Z$ and $\epsilon_\Delta$ differ by an order one coefficient but have the same sign. This leads us to define the efficiency factor $\eta$ as

\begin{align}
Y_{B-L} = \eta(\epsilon_\Delta + 2\epsilon_Z) \Sigma_{\Delta}^\text{eq}_*(T \gg M_\Delta),
\end{align}

so that $\eta \approx 4\eta_0/7$ for $K_{H_u} \ll 1$ (respectively $\eta \approx 7\eta_0/10$ for $K_{H_u} \gtrsim 1$). Comparing the above equations with Eqs. (40) and (41) shows an enhancement of $Y_{B-L}$ by a factor $|(\epsilon_\Delta + 2\epsilon_Z)/\epsilon_\Delta|$, which may be compensated by a smaller $\eta_0$, since $\Delta_{L_{1i}}$ is washed out by the inverse decays of both $\Delta_s$ and $Z_s$. Therefore, we do not expect the final value of the baryon asymmetry to be significantly affected by the presence of the $(\Sigma, \Sigma)$ and $(Z, \overline{Z})$ fields.
4.4 Discussion

The conditions for an order one efficiency,

\[ K_{L_1^c} \ll 1 \, , \, \, K_L, K_{H_u} \gtrsim 1 \, \text{and} \, M_{24} \gg M_\Delta , \]  

have some impact on the value of the CP asymmetry, which is not allowed to be maximal anymore. First, the \( K_{L_1^c} \ll 1 \) condition sets an upper bound on the absolute value of the \( m_{11} \) factor in Eq. (24), which also enters the neutrinoless double beta decay rate. Such a bound can only be satisfied for a normal hierarchy of the light neutrino mass spectrum. Furthermore, the condition \( K_{H_u} \gtrsim 1 \) prevents us from taking a large value of \( \lambda_L \), since Eq. (5) implies \( \lambda_L^2 \approx 0.05 K^{-1}_{H_u} (M_\Delta/10^{12} \text{GeV})(\tau/0.05 \text{eV})^2 \). Finally, avoiding the washout from \((S, T)\)-related processes requires \( M_\Delta \ll M_{24} \), which prevents the loop function from taking its maximal value. As a result, \( \epsilon_\Delta \) takes rather small values in the region of parameter space where the efficiency is of order one, as we discuss below in greater detail.

Let us first recall that, within the uncertainties discussed in Section 3.3, we can take the flavour-blind quantities \( \lambda_L, \lambda_{H_u}, \lambda_{H_d} \) and \( M_\Delta/M_{24} \) as independent parameters (in addition to \( \tan \beta \), the light neutrino masses and the PMNS matrix, which in principle can all have an independent experimental determination). For definiteness, we consider as before the small \( \alpha_d \) case, so that we can in practice set \( \lambda_{H_d} = 0 \), and we choose \( \tan \beta = 10 \) (these parameters do not play a crucial role). We also set \( M_\Delta/M_{24} = 0.1 \), and observe that \( \epsilon_\Delta \) scales linearly with \( M_\Delta/M_{24} \) for the necessarily small values of this ratio. We are then left with \( \lambda_L, \lambda_H \equiv \lambda_{H_u} \) and the light neutrino parameters.

Contour lines of constant \( \epsilon_\Delta \) in the \( \lambda_L-\lambda_H \) plane are shown in Fig. 4 for the case of a normal hierarchical spectrum with \( m_1 \ll m_2 \). In this limit, the asymmetry scales as \( \sin^2 \theta_{13} \) and \( \sin 2\sigma \), and is therefore maximized by \( \sin^2 \theta_{13} = \sin^2 \theta_{13}^{\text{max}} = 0.05 \) and \( \sin 2\sigma = 1 \). Contour lines of constant \( M_\Delta \) are also shown. The efficiency is expected to be large in the unshaded region. In the shaded regions closer to the axes (light blue and red), one of the decay channels \( \Delta_a \to H_u H_u \) or \( \Delta_a \to L\bar{L} \) is out of equilibrium. In the larger shaded region near the \( \lambda_L \) axis (light yellow), the decay channel \( \Delta_a \to L\bar{L} \) is in thermal equilibrium. Fig. 4 shows that the observed value of the baryon asymmetry, \( \eta \epsilon_\Delta \approx 10^{-8} \), can be achieved in a sizeable portion of the parameter space. Although the CP asymmetry grows for growing values of \( m_1 \), the region in which the decay channel \( \Delta_a \to L\bar{L} \) is in thermal equilibrium (implying a small efficiency) becomes simultaneously larger. In the case of an inverted mass hierarchy, the out-of-equilibrium condition is hardly satisfied because \( |m_{11}| \) is bounded from below.

For completeness, we give an example of parameter choice with \( m_1 \sim m_2 \) in the large efficiency region: \( m_1 = 0.005 \text{ eV}, \sin^2 \theta_{15} = 0.05, (\rho, \sigma) = (\pi/4, \pi/2), \lambda_L = 0.1 \gg \lambda_{H_d} \) and \( M_\Delta = 10^{12} \text{ GeV} \). This gives \( \epsilon_\Delta \approx 0.8 \times 10^{-6}(M_\Delta/M_{24}) \) together with \( K_L = 45, K_{H_u} = 5.0 \) and \( K_{L_1} = 0.19 \) (corresponding to \( \lambda_{H_u} = 3.3 \times 10^{-2} \) and \( \lambda_{L_1} = 6.5 \times 10^{-3} \)). We note in passing that this choice of parameters corresponds to an effective Majorana mass \( |m_{ee}| = 3.3 \text{ meV} \) for neutrinoless double beta decay.

A comment is in order on the scale of leptogenesis. As can be seen from Fig. 4, successful leptogenesis in the large efficiency region requires \( M_\Delta \gtrsim 10^{12} \text{ GeV} \). Such values are in strong conflict with the so-called gravitino constraint, which puts an upper bound on the reheating temperature after inflation. From the requirement that
Figure 4: Contours of constant $\epsilon_\Delta$ and $M_\Delta$ in the $\lambda_L$–$\lambda_H$ plane for a normal hierarchical neutrino mass spectrum with $m_1 \ll m_2$, $\sin^2 \theta_{13} = 0.05$ and $\sin 2\sigma = 1$. In the shaded regions, the conditions for a large efficiency are not satisfied.

the gravitino relic density does not exceed the dark matter density, one obtains $T_{RH} \lesssim 10^{9\text{--}10}$ GeV for a gravitino of mass $m_{3/2} \sim 100$ GeV \cite{17}. If the gravitino is not the LSP, a much stronger bound comes from the requirement that its decays do not spoil the successful predictions of Big-Bang Nucleosynthesis \cite{18}, but it is more model-dependent and can be evaded in some schemes (e.g. in the presence of $R$-parity violation \cite{19}). However, there are ways to reconcile the large efficiency regime of our scenario with the gravitino constraint. A first possibility is to assume an extremely light gravitino \cite{20}, $m_{3/2} \leq 16$ eV \cite{21}, where the upper bound comes from WMAP and Lyman-α forest data. In this case, the gravitino decouples when it is still relativistic and escapes the overproduction problem. Such a small gravitino mass can be obtained in some scenarios of gauge-mediated supersymmetry breaking \cite{22}. A second possibility is to assume a very heavy gravitino \cite{23}, $m_{3/2} \gg 100$ TeV. In this case, the gravitino decays before nucleosynthesis and does not affect the light element abundances; furthermore, the (neutralino) LSPs produced in the decays of such heavy gravitinos annihilate efficiently enough to reach their thermal abundance \cite{24}. Therefore, there is no gravitino problem. Finally, a third possibility is to resort to non-thermal production of the heavy triplets, e.g. during reheating \cite{25}, preheating \cite{26}, or via another mechanism \cite{27}. In this way the triplets could be sufficiently produced even though the reheating temperature lies several orders of magnitude below their mass.

Let us note in passing that the level of predictivity of our leptogenesis scenario is maintained in a non-supersymmetric version of the $SO(10)$ model with a real 54 Higgs
multiplet. In this case, there is of course no gravitino problem, but the advantages of supersymmetric unification are lost.

5 Conclusions

We presented a new leptogenesis scenario in which the generated baryon asymmetry depends on the low-energy neutrino parameters. This scenario arises in $SO(10)$ models with Standard Model fermions split among $16$ and $10$ representations and type II realization of the seesaw mechanism. The predictivity of our scenario is due to the fact that the light neutrino masses and the CP asymmetry in triplet decays, $\epsilon_\Delta$, are controlled by the same set of couplings. This is to be contrasted with most leptogenesis models, in which the prediction for the CP asymmetry depends on unknown high-energy flavour parameters. In our scenario, instead, $\epsilon_\Delta$ is proportional to

$$\sum_{k,l} C_{kl} \theta(M_\Delta - M_k - M_l) \text{Im}[m_{kl}(m^* m m^*)_{kl}],$$

where $m$ is the light neutrino mass matrix, and the coefficients $C_{kl}$ depend on the masses of the heavy lepton fields, $M_k$ ($k = 1, 2, 3$). The latter are in turn proportional to the Standard Model charged lepton masses, up to some degree of model dependence. As a result, the CP asymmetry in triplet decays is directly related to the light neutrino parameters; in particular, the CP violation needed for leptogenesis is provided by the CP-violating phases of the PMNS matrix. In the case where $2M_1 < M_\Delta < M_1 + M_2$, $\epsilon_\Delta$ is directly related to the effective Majorana mass of neutrinoless double beta decay.

We discussed the possibility of generating the observed baryon asymmetry in two complementary regimes, assuming for definiteness $M_1 \ll M_\Delta < M_1 + M_2$. The first regime is characterized by large values of the CP asymmetry and by a strong washout. As can be seen in Fig. 3, values of the CP asymmetry as large as $10^{-4} - 10^{-3}$ can be reached for $\theta_{13}$ close to its present upper bound and either a normal mass hierarchy with $m_1 \approx m_2$, or an inverted mass hierarchy with $m_3 \approx \text{few} \times 10^{-2}$ eV. This allows to generate the observed baryon asymmetry for efficiencies as small as $10^{-5} - 10^{-4}$.

A study of the efficiency in this regime requires a detailed numerical analysis. On the contrary, the qualitative features of the large efficiency regime can be studied analytically. Such a regime takes place in a sizeable region of the parameter space where $\lambda_{L,1} \ll \lambda_{H,u}, \lambda_L$ and $M_{24} \gg M_\Delta$. In this region, $\epsilon_\Delta$ takes smaller values, as can be seen in Fig. 4 but still large enough to allow leptogenesis to be successful provided that $\theta_{13}$ is large and the light neutrinos have a normal hierarchical mass spectrum. Upcoming neutrino experiments will further constrain the light neutrino parameters and make it possible to test the viability of the scenario.

Let us finally comment on other low-energy implications of the class of $SO(10)$ models in which our leptogenesis scenario takes place. These models contain heavy states which contribute to the renormalization of the squark and slepton soft supersymmetry breaking masses between high and low energies. The pattern of radiative corrections is not the same as in standard supersymmetric $SO(10)$ models and may lead to distinctive signatures in flavour and CP violating processes. Also, the non-standard

---

9If the $54$ scalar multiplet were complex, the Lagrangian would admit two independent couplings of the $54$ to the $10$, fermion multiplets, $10,10_j (f_{ij} 54^* + g_{ij} 54^*)$, and the connection between the generated baryon asymmetry and the light neutrino parameters would be lost.

---

18
assignment of matter fields in $SO(10)$ representations may affect the predictions for proton decay. We defer the exploration of these effects to future work.

Acknowledgments

We thank T. Hambye, H. Murayama and G. Senjanovic for useful discussions. MF, SL and PH were supported in part by the RTN European Program MRTN-CT-2004-503369 and by the French Program Jeunes Chercheuses et Jeunes Chercheurs of the Agence Nationale de la Recherche (ANR-05-JCJC-0023). MF was supported in part by the Marie Curie Intra-European Fellowship MEIF-CT-2007-039968. PH was supported in part by the Marie Curie Excellence Grant MEXT-CT-2004-01429717.

References

[1] P. Minkowski, Phys. Lett. B 67 (1977) 421; M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, P. van Nieuwenhuizen and D.Z. Freedman (eds.), North Holland Publ. Co., 1979, p. 315; T. Yanagida, in Proc. of the Workshop on the Baryon Number of the Universe and Unified Theories, O. Sawada and A. Sugamoto (eds.), Tsukuba, Japan, 13-14 Feb. 1979, p. 95; S. L. Glashow, in Quarks and Leptons, Cargèse Lectures, 9-29 July 1979, Plenum, New York, 1980, p. 687; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912.

[2] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B 155, 36 (1985).

[3] M. Fukugita and T. Yanagida, Phys. Lett. B 174 (1986) 45.

[4] M. Magg and C. Wetterich, Phys. Lett. B 94 (1980) 61; G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. B 181 (1981) 287; R. N. Mohapatra and G. Senjanovic, Phys. Rev. D 23 (1981) 165. See also J. Schechter and J. W. F. Valle, Phys. Rev. D 22 (1980) 2227.

[5] G. C. Branco, T. Morozumi, B. M. Nobre and M. N. Rebelo, Nucl. Phys. B 617 (2001) 475 [arXiv:hep-ph/0107164]; S. Davidson, J. Garayoa, F. Palorini and N. Rius, Phys. Rev. Lett. 99 (2007) 161801 [arXiv:0705.1503 [hep-ph]].

[6] H. Georgi, AIP Conf. Proc. 23 (1975) 575; H. Fritzsch and P. Minkowski, Annals Phys. 93 (1975) 193.

[7] See e.g. P. H. Frampton, S. L. Glashow and T. Yanagida, Phys. Lett. B 548 (2002) 119 [arXiv:hep-ph/0208157]; G. C. Branco, R. Gonzalez Felipe, F. R. Joaquim, I. Masina, M. N. Rebelo and C. A. Savoy, Phys. Rev. D 67 (2003) 073025 [arXiv:hep-ph/0211001]; S. F. King, Phys. Rev. D 67 (2003) 113010 [arXiv:hep-ph/0211228]; S. Pascoli, S. T. Petcov and W. Rodejohann, Phys. Rev. D 68 (2003) 093007 [arXiv:hep-ph/0302054]; K. Bhattacharya, N. Sahni, U. Sarkar and S. K. Singh, Phys. Rev. D 74 (2006) 093001 [arXiv:hep-ph/0607272]; G. C. Branco, R. Gonzalez Felipe and F. R. Joaquim, Phys. Lett. B 645 (2007) 432 [arXiv:hep-ph/0609297]; S. Pascoli, S. T. Petcov and A. Riotto, Nucl. Phys. B 774 (2007) 1 [arXiv:hep-ph/0611338].
See e.g. M. Plumacher, Nucl. Phys. B 530 (1998) 207 [arXiv:hep-ph/9704231]; E. Nezri and J. Orloff, JHEP 0304 (2003) 020 [arXiv:hep-ph/0004227]; F. Buccella, D. Falcone and F. Tramontano, Phys. Lett. B 524 (2002) 241 [arXiv:hep-ph/0108172]; W. Buchmuller and D. Wyler, Phys. Lett. B 521 (2001) 291 [arXiv:hep-ph/0108216]; G. C. Branco, R. Gonzalez Felipe, E. R. Joaquim and M. N. Rebelo, Nucl. Phys. B 640 (2002) 202 [arXiv:hep-ph/0202030]; J. C. Pati, Phys. Rev. D 68 (2003) 072002 [arXiv:hep-ph/0209160]; E. K. Akhmedov, M. Frigerio and A. Y. Smirnov, JHEP 0309 (2003) 021 [arXiv:hep-ph/0305322]; C. H. Albright and S. M. Barr, Phys. Rev. D 70 (2004) 033013 [arXiv:hep-ph/0404095]; P. Hosteins, S. Lavignac and C. A. Savoy, Nucl. Phys. B 755 (2006) 137 [arXiv:hep-ph/0606078]; E. K. Akhmedov, M. Blennow, T. Hallgren, T. Konstandin and T. Ohlsson, JHEP 0704 (2007) 022 [arXiv:hep-ph/0612194]; S. K. Majee, M. K. Parida, A. Raychaudhuri and U. Sarkar, Phys. Rev. D 75 (2007) 075003 [arXiv:hep-ph/0701109]; J. C. Romao, M. A. Tortola, M. Hirsch and J. W. F. Valle, arXiv:0707.2942 [hep-ph].

[9] J. Hisano, H. Murayama and T. Yanagida, Phys. Rev. D 49 (1994) 4966; Y. Nomura and T. Yanagida, Phys. Rev. D 59 (1999) 017303 [arXiv:hep-ph/9807325]; J. L. Rosner, Phys. Rev. D 61 (2000) 097303; T. Asaka, Phys. Lett. B 562 (2003) 291 [arXiv:hep-ph/0304124]; K. S. Babu, I. Gogoladze, P. Nath and R. M. Syed, Phys. Rev. D 74 (2006) 075004 [arXiv:hep-ph/0607244]; S. M. Barr, Phys. Rev. D 76 (2007) 105024 [arXiv:0706.1490 [hep-ph]].

[10] Z. Berezhiani and Z. Tavartkiladze, Phys. Lett. B 409 (1997) 220 [arXiv:hep-ph/9612232]; Z. Berezhiani and A. Rossi, Nucl. Phys. B 594 (2001) 113 [arXiv:hep-ph/0003084]; M. Malinsky, Phys. Rev. D 77 (2008) 055016.

[11] P. J. O’Donnell and U. Sarkar, Phys. Rev. D 49 (1994) 2118 [arXiv:hep-ph/9307279]; E. Ma and U. Sarkar, Phys. Rev. Lett. 80 (1998) 5716 [arXiv:hep-ph/9802445]; T. Hambye and G. Senjanovic, Phys. Lett. B 582 (2004) 73 [arXiv:hep-ph/0307237].

[12] T. Schwetz, AIP Conf. Proc. 981 (2008) 8 [arXiv:0710.5027 [hep-ph]].

[13] T. Hambye, M. Raidal and A. Strumia, Phys. Lett. B 632 (2006) 667 [arXiv:hep-ph/0510008].

[14] T. Hallgren, T. Konstandin and T. Ohlsson, arXiv:0701.2408 [hep-ph].

[15] E. J. Chun and S. Scopel, Phys. Rev. D 75 (2007) 023508 [arXiv:hep-ph/0609259].

[16] G. Hinshaw et al. [WMAP Collaboration], arXiv:0803.0732 [astro-ph].

[17] M. Bolz, A. Brandenburg and W. Buchmuller, Nucl. Phys. B 606 (2001) 518 [Erratum-ibid. B 790 (2008) 336 [arXiv:hep-ph/0012052]]; J. Pradler and F. D. Steffen, Phys. Rev. D 75 (2007) 023509 [arXiv:hep-ph/0608344]; V. S. Rychkov and A. Strumia, Phys. Rev. D 75 (2007) 075011 [arXiv:hep-ph/0701104].

[18] M. Y. Khlopov and A. D. Linde, Phys. Lett. B 138 (1984) 265; J. R. Ellis, J. E. Kim and D. V. Nanopoulos, Phys. Lett. B 145 (1984) 181. For recent work on the BBN constraints, see e.g. M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, arXiv:0804.3745 [hep-ph].
[19] W. Buchmuller, L. Covi, K. Hamaguchi, A. Ibarra and T. Yanagida, JHEP 0703 (2007) 037 [arXiv:hep-ph/0702184].

[20] H. Pagels and J. R. Primack, Phys. Rev. Lett. 48 (1982) 223.

[21] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese and A. Riotto, Phys. Rev. D 71 (2005) 063534 [arXiv:astro-ph/0501562].

[22] See e.g. M. Ibe, K. Tobe and T. Yanagida, Phys. Lett. B 615 (2005) 120 [arXiv:hep-ph/0503098]; M. Ibe, Y. Nakayama and T. T. Yanagida, arXiv:0804.0636 [hep-ph].

[23] S. Weinberg, Phys. Rev. Lett. 48, 1303 (1982).

[24] N. Arkani-Hamed and S. Dimopoulos, JHEP 0506 (2005) 073 [arXiv:hep-th/0405159]; G. F. Giudice and A. Romanino, Nucl. Phys. B 699 (2004) 65 [Erratum-ibid. B 706 (2005) 65] [arXiv:hep-ph/0406088]; N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B 709 (2005) 3 [arXiv:hep-ph/0409232].

[25] D. J. H. Chung, E. W. Kolb and A. Riotto, Phys. Rev. D 60 (1999) 063504 [arXiv:hep-ph/9809453].

[26] G. N. Felder, L. Kofman and A. D. Linde, Phys. Rev. D 59 (1999) 123523 [arXiv:hep-ph/9812289]; G. F. Giudice, M. Peloso, A. Riotto and I. Tkachev, JHEP 9908 (1999) 014 [arXiv:hep-ph/9905242].

[27] G. F. Giudice, A. Riotto and A. Zaffaroni, Nucl. Phys. B 710 (2005) 511 [arXiv:hep-ph/0408155]; R. Allahverdi and A. Mazumdar, JCAP 0610 (2006) 008 [arXiv:hep-ph/0512227]; G. F. Giudice, L. Mether, A. Riotto and F. Riva, arXiv:0804.0166 [hep-ph].