Higher-Dimensional Origin of Heavy Sneutrino Domination and Low-Scale Leptogenesis

A. Hebecker[a], J. March-Russell[a], and T. Yanagida[a,b]

[a] *Theory Division, CERN, CH-1211 Geneva 23, Switzerland*
[b] *Department of Physics, University of Tokyo, Tokyo 113, Japan*

(August 27, 2002)

Abstract

If the expectation value of the right-handed (rhd) sneutrino comes to dominate the universe, its decay naturally leads to successful leptogenesis, as well as significant dilution of dangerous inflationary relics, such as the gravitino. The resulting baryon asymmetry is independent of other cosmological initial conditions. This attractive variant of leptogenesis requires at least one of the rhd neutrinos to have small Yukawa coupling and to have mass $\sim 10^6$ GeV, much smaller than the grand unified (GUT) scale. We show that these features naturally arise in the context of independently motivated and successful 5d orbifold GUTs with inverse-GUT-scale-sized extra dimensions. Rhd neutrinos are realized as bulk fields N_i with 5d bulk masses, while Yukawa couplings and lepton-number-violating masses for the N_i are localized at the SM boundary. The exponential suppression of the would-be N_i zero-modes leads to the desired small 4d Yukawa couplings and small masses for the rhd neutrino states. The see-saw prediction for the lhd neutrino mass scale is automatically maintained. We show that this realization of rhd neutrinos is nicely accommodated within an attractive orbifold-GUT flavour model, where all flavour hierarchies have a geometrical origin.

1 Basics of the scenario

With growing experimental evidence for neutrino masses in a range that is consistent with a GUT-scale-based see-saw mechanism [1], leptogenesis has become the standard scenario for the generation of the baryon asymmetry of the universe. In the original proposal [2], the heavy rhd neutrinos decay in an out-of-equilibrium fashion once the universe has cooled to a temperature below their mass scale. The resulting lepton number is then converted to baryon number by standard model (SM) sphaleron processes (for a recent review see, e.g., [3]). Alternatively, in a supersymmetric theory, lepton number can be generated by the decay of a condensate of the scalar component $\tilde{N}$ of the rhd neutrino superfield N [4]. In particular, a recent detailed analysis [5] of this scenario has shown that, if $\tilde{N}$ comes to dominate the universe, its decay can naturally produce the required lepton asymmetry independently of other cosmological initial conditions. At the same time, the number density of dangerous inflationary relics, such as the gravitino, is significantly diluted.

This cosmologically attractive variant of leptogenesis requires at least one of the rhd neutrino masses to be very small compared with the GUT scale, $\sim 10^6\,\text{GeV}$, and the corresponding Yukawa coupling to be suppressed. In the present paper, we show that such a situation arises naturally in the context of independently motivated higher-dimensional GUTs (with inverse-GUT-scale-sized extra dimensions). Before embarking, in Sects. 2 and 3, upon a detailed discussion of the orbifold-GUT model and of the cosmology, we now explain the fundamentals of our scenario.

To be specific, we will formulate our ideas in the framework of supersymmetric SU(5) orbifold GUTs [6–9] in 5 dimensions. These theories are attractive because they incorporate the success of MSSM gauge-coupling unification [8, 9], while providing natural doublet-triplet splitting as well as suppressed proton decay [6–9]. Moreover, a three generation model with a geometrical origin of hierarchical Yukawa couplings and see-saw neutrinos can easily be realized [10].

The starting point for these models is a 5d super Yang-Mills theory compactified on an interval with coordinate $y \in [0, l]$. At the $y = 0$ boundary (the 'SU(5) brane') the 5d gauge symmetry is unbroken by boundary conditions, while at $y = l$ (the 'SM brane'), the boundary conditions on the gauge fields explicitly break the 5d SU(5) down to the SM gauge group. Below the compactification scale $M_c \equiv 1/l$, one has a 4d effective field theory with SM gauge group (and N=1 supersymmetry).

A basic property of such models is that the 'bulk', $y \in (0, l)$, is moderately large compared to the fundamental 5d Planck (or UV cutoff) length. Arguments pointing to this conclusion include the weakness of the effective 4d unified coupling, the 'observed' smallness of GUT-scale threshold corrections, and the flavour hierarchies among the generations (see e.g. [8, 9, 11]). As we discuss in Sect. 2, an alternative way of quantifying the size of the bulk follows from the requirement that gauge coupling unification (including the KK-mode corrected logarithmic running above M_c) occurs at the 5d Planck scale. Concretely this argument favours an orbifold GUT setup with $Ml \simeq 300$ and $M \simeq 1.4 \times 10^{17}\,\text{GeV}$, where M is the reduced 5d Planck mass. This is in accord with the size of the bulk deduced from other arguments.

If a bulk mass m, odd under 5d parity, is introduced for a bulk matter field, its zero mode develops an exponential profile $\sim \exp[-ym]$ [12] (see also [13,11] and [14]). Thus, depending on the sign of m, zero modes can be strongly peaked at either brane. If one of the three SM-singlet rhd neutrino fields N_i is exponentially peaked at the SU(5) brane, while lepton number violating mass terms $\sim N_i^2$ are allowed only at the SM brane, an exponential suppression of both the 4d rhd neutrino mass and Yukawa coupling naturally arises. The crucial observation is that the light neutrino masses are not affected even if one or more of the N_i are such bulk fields with arbitrary bulk profile. This is clear because the N_i zero modes receive only their kinetic term from the bulk, while their effective 4d mass and Yukawa coupling come from the brane. When the N_i are integrated out, their kinetic term plays no role and thus it is irrelevant whether they are brane or bulk fields. Therefore the traditional see-saw prediction for the lhd neutrino mass scale is maintained. However, the rhd sneutrino mass scale is exponentially suppressed, as are its Yukawa couplings, and thus decay width. These are the new features that allow us to realise the attractive scenario of sneutrino ($\tilde{N}$) dominated cosmology and leptogenesis.

In Sect. 2 we present a more detailed motivation and quantitative analysis of the basic orbifold picture of neutrino masses and interactions, in particular a demonstration that it can be successfully embedded in a full flavour model. Specifically, both Higgs doublets and the three $\overline{\mathbf{5}}$'s of SU(5) (denoted by $\overline{F}_i$) are localized at the SM brane, while the three $\mathbf{10}$'s (denoted by T_i) are bulk fields. The quark and lepton mass hierarchies are generated by the bulk profiles of the T_is.

However, we emphasise that our neutrino mass construction is quite generic and does not rely on the details of the specific SU(5) model worked out in the rest of this paper. The crucial ingredient is a 5d, or higher-dimensional theory compactified on an interval with Yukawa couplings and lepton-number-violating neutrino masses localized at one of the boundaries. The exponential suppression of zero-mode wave functions at that boundary generates both the fermion mass hierarchy and the desired light rhd neutrinos.

The cosmology of the above model of neutrinos has many attractive aspects. In particular, over a wide parameter region it leads to the $\tilde{N}$ dominated early universe of [5], as $\tilde{N}$ has an exponentially enhanced life-time. In more detail, if the initial value of $|\tilde{N}|$ is of the order of M, a natural circumstance, then $\tilde{N}$ will come to dominate the universe for inflationary reheating temperatures $T_R \gtrsim 10^9 \, \text{GeV}$. Moreover, if T_R is varied between $\sim 10^9 \, \text{GeV}$ and $\sim 10^{12} \, \text{GeV}$, the gravitino number density in the late universe remains fixed at the level corresponding to $T_R \sim 10^9 \, \text{GeV}$. This attractive feature of $\tilde{N}$ dominated cosmology is due to the entropy produced by $\tilde{N}$ decay. Finally, the decay of $\tilde{N}$ produces the lepton-number asymmetry.

This cosmology is a fascinating possibility since most of the important physical parameters in the present universe, such as baryon-number asymmetry, entropy (and, as we later discuss, even spectrum of density fluctuations), are determined by the nature of the scalar partner of the lightest rhd neutrino. In Sect. 3, we provide a more detailed discussion of this cosmology, while some further possibilities, together with our conclusions, are contained in Sect. 4.

2 The flavour model

Consistency of the orbifold GUT framework requires $M_c = 1/l$ to be significantly smaller than the UV scale M of the 5d gauge theory. To be more specific, the (reduced) Planck masses in 4d $(\overline{M}_P)$ and in 5d (M), are related by

$$\overline{M}_P^2 = M^3 \, l \,, \qquad\qquad \overline{M}_P = M_P/\sqrt{8\pi} \simeq 2.4 \times 10^{18}\,\mathrm{GeV} \,, \qquad (1)$$

and we demand gauge coupling unification at the fundamental scale M. In spite of the UV sensitivity of the non-renormalizable 5d theory, the differences of inverse SM gauge couplings $\alpha_{ij} = \alpha_i^{-1} - \alpha_j^{-1}$ $(i = 1, 2, 3)$ continue to run logarithmically above M_c [8,9] because these differences are only sensitive to the SU(5)-breaking SM brane. In the context of the minimal model of [9], where the Higgs-doublets are localized at the SM brane, this 'differential running' [15] comes entirely from the gauge sector. With the effective SUSY breaking scale set to m_Z, we have

$$\alpha_{ij}(m_Z) = \alpha_{ij}(M) + \frac{1}{2\pi} \left\{ a_{ij} \ln \frac{M}{m_Z} + \frac{1}{2} b_{ij} \ln \frac{M}{M_c} \right\} \,, \qquad (2)$$

where $a_{ij} = a_i - a_j$ and $b_{ij} = b_i - b_j$ (with $a_i = (33/5, 1, -3)$ and $b_i = (-10, -6, -4)$) characterise the familiar MSSM running and the KK mode contributions respectively. If we define the conventional 4d unification scale by the meeting of the U(1) and SU(2) couplings α_1 and α_2, then the low-energy data $\alpha_i^{-1}(m_Z) = (59.0, 29.6, 8.4)$ imply $M_{\mathrm{GUT}} \simeq 1.9 \times 10^{16}\,\mathrm{GeV}$. By contrast, combining Eqs. (1) and (2) and assuming $\alpha_{12}^{-1}(M) = 0$, one derives the 5d unification scale $M = 1.4 \times 10^{17}\,\mathrm{GeV}$ and $Ml = 2.8 \times 10^2$. Of course, these numbers represent only rough estimates since the $\alpha_{ij}(M)$ have, in general, non-zero $\mathcal{O}(1)$ values, which perturb the calculation of Ml and M. (In the slightly different approach of [16], the model is fixed by requiring the precision of simultaneous 1-2 and 2-3 unification to be better than with conventional MSSM running.)

The above discussion provides us with a motivation for an orbifold GUT setup with $M \simeq 1.4 \times 10^{17}\,\mathrm{GeV}$ and with the small parameter $\varepsilon^2 \equiv 1/(Ml) \simeq 1/300$. Flavour is described by introducing the three $\overline{F}_i$ fields and the two Higgs doublets H_u and H_d on the SM brane (recall that, because of the reduced symmetry of the SM brane, there is no need for Higgs triplets), while allowing the T_i to propagate in the bulk.[1] This large disparity between T's and $\overline{F}$'s is the geometric origin of small quark and large lepton mixing (cf. [17] and [10]). If a bulk mass m, odd under 5d parity, is introduced for a 5d hypermultiplet, its zero mode develops an exponential profile $\sim \exp[-ym]$. Thus, depending on the sign of m, zero modes can be strongly peaked at either brane. In particular, this allows for a dynamical realization of SM-brane fields with quantum numbers appropriate for an SU(5) representation (e.g., the Higgs doublets and the $\overline{F}_i$ above). We will use this additional tool to realize the fermion mass hierarchy by appropriately localizing the T_i.

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[1]To be more precise, one introduces hypermultiplets T_i and T_i' and assigns boundary conditions ensuring that the zero modes correspond to the field content of three **10**'s of SU(5) [8,9]. Moreover, the boundary conditions break N=2 to N=1 SUSY, leaving us with conventional chiral multiplets at low energy $E < M_c$.

An understanding of the observed hierarchies in the fermion masses and mixings emerges naturally if the bulk mass of T_3 is sufficiently large and negative, $m < 0$ (so that, for all practical purposes, T_3 is a SM brane field), while T_2 has vanishing bulk mass (flat zero-mode) and T_1 has a finite bulk mass $m > 0$ (its zero-mode therefore being suppressed at the SM brane).

Concretely, the primordial and unstructured $\mathcal{O}(1)$ Yukawa couplings λ at the SM brane are rescaled as

$$\lambda \to \frac{\lambda}{\sqrt{Ml}} \sqrt{\frac{2ml}{e^{2ml} - 1}} \tag{3}$$

for each participating bulk field with bulk mass m. This rescaling follows from the 4d canonical normalization of the 5d kinetic term and the exponentially suppressed zero-mode field value at the SM brane. Applying this to T_2, one finds that this field enters Yukawa interactions with a suppression factor ε. The analogous suppression factor for T_1 depends on m and becomes $\sim \varepsilon^2$ for the choice $ml \simeq 3.9$. This leads to the following realistic Yukawa matrix structure for the two effective 4d interactions $H_u T^T \lambda_{TT} T$ and $H_d T^T \lambda_{TF} \overline{F}$:

$$\lambda_{TT} \sim \begin{pmatrix} \varepsilon^4 & \varepsilon^3 & \varepsilon^2 \\ \varepsilon^3 & \varepsilon^2 & \varepsilon \\ \varepsilon^2 & \varepsilon & 1 \end{pmatrix}, \qquad \lambda_{TF} \sim \begin{pmatrix} \varepsilon^2 & \varepsilon^2 & \varepsilon^2 \\ \varepsilon & \varepsilon & \varepsilon \\ 1 & 1 & 1 \end{pmatrix}, \tag{4}$$

with unknown $\mathcal{O}(1)$ factors multiplying each entry.[2] It is known that this Yukawa coupling hierarchy also gives rise to an approximately correct CKM structure. The top Yukawa coupling is naturally $\mathcal{O}(1)$. The required relative suppression of down-type Yukawa couplings can be realized either by going to large $\tan\beta$ or by slightly decreasing the strength with which H_d is peaked at the SM brane.

The construction presented so far can be summarized as follows. By identifying the 5d Planck mass with the unification scale, we have argued for a relative bulk size characterized by $\varepsilon \sim 1/\sqrt{Ml} \sim 1/\sqrt{300}$. If all fields except T_1 and T_2 are localizing at the SM brane, this bulk suppression factor beautifully explains the mass hierarchy between the two heavier generations [11]. To explain the extreme lightness of the first generation, we had to give T_1 a bulk profile exponentially suppressed at the SM brane using the additional tool of bulk masses. With this tool in hand, rhd neutrino singlets can easily acquire the exponentially suppressed 4d masses and couplings required for the $\tilde{N}$ dominated universe.

Now we discuss the rhd neutrinos in more detail. Consider introducing three neutrino fields N_i at the SM brane. Given a Majorana mass matrix $M_{N,ij}$ with $\mathcal{O}(M)$ entries and $\mathcal{O}(1)$ Yukawa couplings between N_i, $\overline{F}_i$ and H_u, the conventional see-saw mechanism leads to a light neutrino mass scale $|H_u|^2/M \simeq 2 \times 10^{-4}$ eV. In the present scenario, such a small mass scale is welcome since it ensures the out-of-equilibrium decay of $\tilde{N}$ (see Sect. 3). The observed neutrino oscillations, which require a somewhat larger

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[2]A slight modification, leading to a welcome further suppression of electron and down-quark mass, is obtained by placing one of the $\overline{F}$'s in the bulk. In fact, such a construction can be motivated by its particularly high symmetry: One set of fields $(T, \overline{F})$ are on the SM brane, one set are massless bulk fields, and the third set are massive bulk fields with the sign of the mass flipped between T and $\overline{F}$ (making $\overline{F}$ effectively a SM brane field).

light neutrino mass scale, can be accommodated by assuming that M_N has two slightly suppressed eigenvalues. (A concrete example of such a suppression mechanism will be provided shortly.)

As discussed in Sect. 1, light neutrino masses are not affected if one or more of the N_i are promoted to bulk fields. To be specific, let us declare N_1 to be a bulk field with bulk mass m_1 and effective 4d mass (cf. Sect. 5 of [10])

$$M_1 \simeq 2m_1 e^{-2m_1 l} \, . \tag{5}$$

Due to the exponential suppression factor, the desired small value of M_1 is easily realized: for example, $M_1 \simeq 3 \times 10^6 \, \text{GeV}$ for $m_1 l \simeq 11$. While this concludes the description of our basic flavour model with a naturally light rhd neutrino, several open issues deserve further discussion.

Firstly, we need to enhance two of the light neutrino masses. For example, one could introduce a Froggatt-Nielsen U(1), broken by two charge-(± 1) fields with vacuum expectation values $|\Phi_\pm|$ where $|\Phi_\pm|/M \simeq \eta \ll 1$ [18]. With charge assignments $(0,-1,-1)$ and $(1,1,1)$ for the N_i and $\overline{F}_i$ respectively, one obtains the following structures for the Yukawa matrix λ_N in $H_u \overline{F}_i^T \lambda_{N,ij} N_j$ and the mass matrix:

$$\lambda_N \sim \begin{pmatrix} \eta & 1 & 1 \\ \eta & 1 & 1 \\ \eta & 1 & 1 \end{pmatrix} , \qquad M_N \sim \begin{pmatrix} 1 & \eta & \eta \\ \eta & \eta^2 & \eta^2 \\ \eta & \eta^2 & \eta^2 \end{pmatrix} . \tag{6}$$

It is easy to convince oneself that all entries of the resulting light neutrino mass matrix $m_\nu \simeq \lambda_N M_N^{-1} \lambda_N^T \times |H_u|^2$ are of the order $\eta^{-2}|H_u|^2/M$. This also sets the scale for two of the eigenvalues. Although the remaining eigenvalue is suppressed to $\eta^2|H_u|^2/M$, all three mixing angles are generically large.[3] In our setup, realistic neutrino phenomenology requires $\eta \sim 10^{-1}$. Furthermore, assigning a suitable U(1) charge to H_d provides an alternative way to realize suppressed down-type masses. Let us finally argue why the family-symmetry should be broken in the U(1) charge assignment of the N_i. One possibility is to demand vanishing U(1) charges for all bulk fields. Alternatively, one could replace the U(1) with a Z_3 and then note that, while the cancellation of the mixed Z_3-SU(5) anomaly forces all three $\overline{F}_i$ to have the same charge, the N_i charges remain unrestricted.

Secondly, it is necessary to forbid both parity-even bulk masses as well as SU(5)-brane-localized mass terms for the rhd neutrinos. Following [10], this can be done by gauging $U(1)_\chi$ (named as in [20]), defined by $SU(5) \times U(1)_\chi \subset SO(10)$. Since $\tilde{N}$ domination requires a large initial value of $\tilde{N}_1$, the D-term potential for $\tilde{N}_1$ has to be suppressed. This can be achieved by dynamically breaking $U(1)_\chi$ at the high scale M in 5d. A surviving discrete subgroup will be sufficient to forbid the dangerous mass operators. In addition, it is natural that $U(1)_\chi$ is broken by orbifolding at the SM brane [10], making it the only possible location for the required lepton-number violating mass term.

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[3]The smallness of 1-3 mixing may be accidental (cf. [19]). Alternatively, it could be explained in a modified model where one of the $\overline{F}$'s is a bulk field.

Furthermore, we would like to comment on the relation of our method of generating the fermion mass hierarchy and the light N_1 field to the familiar Froggatt-Nielsen approach. Certainly the assignment of bulk masses to different sets of fields resembles the assignment of U(1) charges. This similarity becomes even more pronounced if the bulk masses are dynamically realized by expectation values of U(1) fields with Fayet-Iliopoulos terms at the boundary (see, e.g., [21]). However, especially in the case of the large suppression factor needed for N_1, it is a significant advantage that the bulk mass effect is exponential rather than power-like. Furthermore, there are crucial qualitative differences in the resulting phenomenology. For example, higher-order Kähler-terms involving $T_i^\dagger T_i$ together with the SUSY-breaking spurion, which can lead to dangerous flavour violation, are unrestricted by U(1) symmetries. In our case, if SUSY breaking is localized at the SM brane, such terms will be geometrically suppressed for the first two generations. The argument extends to the $\overline{F}_i$ if some of them are promoted to bulk fields.

Finally, note that the bulk masses used in the above construction are significantly smaller than the fundamental 5d scale M. This may follow naturally if bulk masses come from expectation values of weakly coupled U(1) fields. Alternatively, one may imagine the 5d theory to descend from a 6d theory, where bulk masses are forbidden, so that 5d bulk masses are due to small, non-perturbative effects arising in the 6d to 5d compactification process.

3   Cosmology

Let us turn to the discussion of cosmology. It is a reasonable assumption that the scalar partner of at least one of the right-handed neutrinos has, during inflation, an expectation value of the order of the cutoff scale M. The reason for this is that higher-dimension operators in the Kähler potential link the inflationary sector, in particular the superfield whose F-term or D-term gives rise to the non-zero vacuum energy, and the rhd neutrino superfields. This leads to a contribution to the (mass)^2 of the sneutrino of order $(H_{\text{inf}})^2$, where H_{inf} is the expansion rate of inflation, with sign that depends upon the unknown Kähler operator coefficient. If this sign is such that the (mass)^2 is negative, then $\tilde{N}_1$ gains an expectation value only limited by yet higher-order terms in the potential, whose natural scale is M. Here we have assumed that H_{inf} is larger than the mass of the right-handed neutrino.

As described in the previous sections, for the lightest rhd sneutrino field $\tilde{N}_1$ both its mass M_1 (cf. Eq. 5) and its effective 4d Yukawa coupling,

$$\lambda_{N,i1} \simeq \eta\sqrt{2m_1/M}\,e^{-m_1 l}\,, \tag{7}$$

are exponentially suppressed. When the expansion rate H after the end of inflation decreases below M_1, $\tilde{N}_1$ starts coherently oscillating. Given a condition on the post-inflationary reheating temperature T_R (to be discussed below), the oscillation energy dominates the energy density of the early universe, and its decay produces the baryon asymmetry observed today without any cosmological difficulty.

If the Yukawa couplings $\lambda_{N,ij}$ of the N_j have CP violating phases, the decay of $\tilde{N}_1$ produces a lepton-number asymmetry ε_1 given by [22,5]

$$\varepsilon_1 \simeq 1 \times 10^{-10}(M_1/10^6\,\text{GeV})(m_{\nu_3}/0.05\,\text{eV})\delta_{eff}. \qquad (8)$$

Here, δ_{eff} is an effective CP violating phase. This lepton asymmetry is converted into a combined baryon and lepton asymmetry through non-perturbative electroweak sphaleron effects. A crucial observation of ref. [5] is that the final baryon asymmetry is determined by the reheating temperature, T_{N_1}, of the $\tilde{N}_1$ decay once it dominates the energy density of the early universe. The net baryon to entropy ratio is given by [5]

$$n_B/s \simeq (8/23)(3/4)\varepsilon_1(T_{N_1}/M_1) \simeq 0.3 \times 10^{-10}(T_{N_1}/10^6\,\text{GeV})(m_{\nu_3}/0.05\,\text{eV})\delta_{eff}, \qquad (9)$$

where n_B and s are baryon-number and entropy densities, respectively. The observed baryon asymmetry $n_B/s \simeq (0.4-1) \times 10^{-10}$ is obtained by taking $T_{N_1} \simeq 2 \times 10^6\,\text{GeV}$ for $\delta_{eff} \simeq 1$. On the other hand, the reheating temperature T_{N_1} due to $\tilde{N}_1$ decay is given by $T_{N_1}^2 \simeq \Gamma_{N_1}\overline{M}_P$, where the decay rate Γ_{N_1} is

$$\Gamma_{N_1} \simeq (3/4\pi)\lambda_{N,i1}^2 M_1. \qquad (10)$$

We see that the desired reheating temperature $T_{N_1} \simeq 2 \times 10^6\,\text{GeV}$ is obtained for $m_1 l = 11$, where we have used Eqs. (5) and (7). Notice that the rhd neutrino mass $M_1 \simeq 3 \times 10^6\,\text{GeV} > T_{N_1}$, and hence the out-of-equilibrium condition for $\tilde{N}_1$ decay is automatically satisfied.

Let us now discuss the condition for the $\tilde{N}_1$ domination in the early universe. Since the initial value of $|\tilde{N}_1| \simeq 10^{17}\,\text{GeV}$, the energy density of the coherent $\tilde{N}_1$ oscillation, at the start of this oscillation, is a minor component of the total density. However, if the $\tilde{N}_1$ lifetime is sufficiently longer than that of the inflaton, it can dominate the early universe since the energy density of the radiation resulting from the inflaton decay dilutes faster than the energy density of the coherent oscillation. Thus, the condition for the $\tilde{N}_1$ domination is translated to an upper limit on the inflaton lifetime for a given $\tilde{N}_1$ lifetime. Written in terms of the post-inflationary reheating temperature, T_R, this condition is

$$T_R > 3T_{N_1}(\overline{M}_P/|\tilde{N}_1|)^2 \simeq 2 \times 10^9\,\text{GeV}, \qquad (11)$$

which is easily satisfied in a variety of inflationary models.

The post-inflationary reheating temperature must also satisfy an upper bound so as to avoid the over-production of gravitinos. In the standard cosmology (without $\tilde{N}_1$ domination) the upper bound is determined to be $T_R < 10^{10}\,\text{GeV}$ for a gravitino mass $\sim 1\,\text{TeV}$ [23] (for a recent analysis see also [24]), giving a stringent restriction on inflationary models.

However, with $\tilde{N}$-dominated cosmology, only a weaker constraint applies. The reason for this is that $\tilde{N}_1$ decay reheats the universe once more, and the associated entropy production dilutes substantially the density of earlier-produced gravitinos. In detail, start from a situation where T_R is at its lower bound given by Eq. (11) and raise the reheating temperature gradually. While the post-inflationary gravitino production increases proportionally to T_R, the subsequent $\tilde{N}_1$ decay introduces the dilution factor

$$1/\Delta \simeq 3T_{N_1}/T_R(\overline{M}_P/|\tilde{N}_1|)^2, \qquad (12)$$

which precisely compensates the previous effect. Thus, the number density of gravitinos is determined by an effective reheating temperature $T_{R,eff}$ instead of the original T_R. This effective temperature is given by

$$T_{R,eff} = 1/\Delta \times T_R = 3T_{N_1}(\overline{M}_P/|\tilde{N}_1|)^2 \simeq 2 \times 10^9 \,\text{GeV}. \tag{13}$$

However, if the reheating temperature rises above $T_R \simeq 10^{12}$ GeV, the situation changes because now reheating takes place before $\tilde{N}_1$ oscillations start. While the initial gravitino production continues to grow with T_R, the dilution factor remains constant, giving

$$T_{R,eff} = 1/\Delta \times T_R = 3T_{N_1}(\overline{M}_P/|\tilde{N}_1|)^2(T_R/10^{12}\,\text{GeV}) \simeq 2 \times 10^9 \,\text{GeV}(T_R/10^{12}\,\text{GeV}). \tag{14}$$

Applying the analysis of Ref. [23], we see that a gravitino of mass of order $1\,\text{TeV}$ is consistent with a significantly extended range of post-inflationary reheating temperatures, $T_R < 10^{13}\,\text{GeV}$. Putting this upper bound together with our earlier lower bound from $\tilde{N}$ domination leads to an allowed range, $2 \times 10^9\,\text{GeV} < T_R < 10^{13}\,\text{GeV}$, for successful sneutrino-dominated leptogenesis.

4   Conclusions

In summary, we have presented a higher-dimensional scenario, well motivated from a particle-physics perspective, in which cosmological heavy sneutrino domination occurs naturally and low-scale leptogenesis is responsible for the observed baryon asymmetry of the universe. The entropy produced in the decay of the $\tilde{N}$ condensate dilutes unwanted relics from the period of reheating, alleviating in particular the danger of gravitino overproduction.

We briefly comment on density perturbations in this scenario. If the dominant density perturbation originates from the fluctuations of $\tilde{N}$ during inflation, the scale-invariant spectrum is automatic and observations do not give stringent constraints on the model of inflation [25]. Since, in our scenario, $\tilde{N}$ decay is the origin of baryon number asymmetry and dark matter, we necessarily have adiabatic perturbation dominance [26].

Thus, the role of the inflaton is reduced to providing a period of exponential expansion while its two main dynamical effects, the production of density perturbations and the reheating of the universe, are taken over by the heavy sneutrino. In fact, the presence of the inflaton mainly has a constraining effect – it has to decay sufficiently early to allow for heavy sneutrino domination and sufficiently late not to produce an excess of gravitinos. Thus, one might also wonder whether it is possible to get rid of the inflaton altogether. One obvious possibility would be to assume a sneutrino potential with a flat region away from the origin, so that a sufficiently long inflationary period driven by the sneutrino condensate is realized. It would be a very interesting and challenging task to understand the origin of such an unusual sneutrino potential.

Finally, we discuss a less radical way of avoiding the constraints associated with the decay of the inflaton entirely. If the inflaton potential is such that, in the true vacuum, the

inflaton is massless, its energy density during the oscillation period decays faster[4] than that of $\tilde{N}$, which varies with the scale factor R as R^{-3}. Thus, the desired $\tilde{N}$ domination is always obtained. The masslessness of the inflaton is technically natural since we do not require any non-gravitational coupling of the inflaton to the matter sector. It would be interesting to write down and analyse a well-motivated and complete inflation model with a potential that leads to such a 'harmless' late time behaviour of the inflaton.

References

[1] T. Yanagida, in Proceedings of the *Workshop on Unified Theories and Baryon Number of the Universe*, editors; O.Sawada and A.Sugamoto, (Tsukuba, Japan, 1979, KEK Report KEK-79-18) p. 95;
M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity* (North Holland, Amsterdam, 1979) eds. P. van Nieuwenhuizen and D. Freedman, p. 315.

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

[3] W. Buchmüller and M. Plümacher, Int. J. Mod. Phys. A **15** (2000) 5047 [arXiv:hep-ph/0007176].

[4] H. Murayama and T. Yanagida, Phys. Lett. B **322** (1994) 349 [arXiv:hep-ph/9310297].

[5] K. Hamaguchi, H. Murayama and T. Yanagida, Phys. Rev. D **65** (2002) 043512 [arXiv:hep-ph/0109030].

[6] Y. Kawamura, Prog. Theor. Phys. **105** (2001) 999 [arXiv:hep-ph/0012125].

[7] G. Altarelli and F. Feruglio, Phys. Lett. B **511** (2001) 257 [arXiv:hep-ph/0102301].

[8] L. J. Hall and Y. Nomura, Phys. Rev. D **64** (2001) 055003 [arXiv:hep-ph/0103125].

[9] A. Hebecker and J. March-Russell, Nucl. Phys. B **613** (2001) 3 [arXiv:hep-ph/0106166]; Nucl. Phys. B **625** (2002) 128 [arXiv:hep-ph/0107039].

[10] A. Hebecker and J. March-Russell, Phys. Lett. B **541** (2002) 338 [arXiv:hep-ph/0205143].

[11] L. Hall, J. March-Russell, T. Okui and D. R. Smith, arXiv:hep-ph/0108161; L. J. Hall, Y. Nomura and D. R. Smith, arXiv:hep-ph/0107331.

[12] R. Jackiw and C. Rebbi, Phys. Rev. D **13** (1976) 3398; D. B. Kaplan, Phys. Lett. B **288** (1992) 342 [arXiv:hep-lat/9206013].

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[4]We are grateful to J. Garcia-Bellido for helpful discussions on this point.

[13] N. Arkani-Hamed *et al.*, Phys. Rev. D **65** (2002) 024032 [arXiv:hep-ph/9811448]; N. Arkani-Hamed and S. Dimopoulos, Phys. Rev. D **65** (2002) 052003 [arXiv:hep-ph/9811353].

[14] H. Georgi, A. K. Grant and G. Hailu, Phys. Rev. D **63** (2001) 064027 [arXiv:hep-ph/0007350]; D. E. Kaplan and T. M. Tait, JHEP **0111** (2001) 051 [arXiv:hep-ph/0110126].

[15] Y. Nomura, D. R. Smith and N. Weiner, Nucl. Phys. B **613** (2001) 147 [arXiv:hep-ph/0104041].

[16] L. J. Hall and Y. Nomura, Phys. Rev. D **65** (2002) 125012 [arXiv:hep-ph/0111068].

[17] T. Watari and T. Yanagida, arXiv:hep-ph/0205090.

[18] C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B **147** (1979) 277.

[19] L. J. Hall, H. Murayama and N. Weiner, Phys. Rev. Lett. **84** (2000) 2572 [arXiv:hep-ph/9911341].

[20] D. E. Groom *et al.* [Particle Data Group], Eur. Phys. J. C **15** (2000) 1 (see p. 290).

[21] R. Barbieri *et al.*, arXiv:hep-th/0203039; S. Groot Nibbelink, H. P. Nilles and M. Olechowski, arXiv:hep-th/0205012; D. Marti and A. Pomarol, arXiv:hep-ph/0205034.

[22] W. Buchmüller and T. Yanagida, Phys. Lett. B **445** (1999) 399 [arXiv:hep-ph/9810308].

[23] M. Y. Khlopov and A. D. Linde, Phys. Lett. B138, 265 (1984); J. Ellis, J. E. Kim and D. V. Nanopoulos, Phys. Lett. B145, 181 (1984); M. Kawasaki and T. Moroi, Prog. Theor. Phys. 93, 879 (1995) [arXiv:hep-ph/9403364].

[24] M. Kawasaki, K. Kohri and T. Moroi, Phys. Rev. D63, 103502 (2001) [arXiv:hep-ph/0012279].

[25] D. H. Lyth and D. Wands, Phys. Lett. B **524** (2002) 5 [arXiv:hep-ph/0110002]; T. Moroi and T. Takahashi, Phys. Lett. B **522** (2001) 215 [Erratum-ibid. B **539** (2002) 303] [arXiv:hep-ph/0110096].

[26] T. Moroi and T. Takahashi, arXiv:hep-ph/0206026.