Sneutrino Inflation in the Light of WMAP: Reheating, Leptogenesis and Flavour-Violating Lepton Decays

John Ellis\textsuperscript{1}, Martti Raidal\textsuperscript{1,2} and T. Yanagida\textsuperscript{3}

\textsuperscript{1} TH Division, CERN, CH-1211 Geneva 23, Switzerland
\textsuperscript{2} National Institute of Chemical Physics and Biophysics, Tallinn 10143, Estonia
\textsuperscript{3} Department of Physics, University of Tokyo, Tokyo 113-0033, Japan

ABSTRACT

We reconsider the possibility that inflation was driven by a sneutrino - the scalar supersymmetric partner of a heavy singlet neutrino - in the minimal seesaw model of neutrino masses. We show that this model is consistent with data on the cosmic microwave background (CMB), including those from the WMAP satellite. We derive and implement the CMB constraints on sneutrino properties, calculate reheating and the cosmological baryon asymmetry arising via direct leptogenesis from sneutrino decays following sneutrino inflation, and relate them to light neutrino masses. We show that this scenario is compatible with a low reheating temperature that avoids the gravitino problem, and calculate its predictions for flavour-violating decays of charged leptons. We find that $\mu \rightarrow e\gamma$ should occur close to the present experimental upper limits, as might also $\tau \rightarrow \mu\gamma$. 

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1 Introduction

Inflation \(^1\) has become the paradigm for early cosmology, particularly following the recent spectacular CMB data from the WMAP satellite \(^2\), which strengthen the case made for inflation by earlier data, by measuring an almost scale-free spectrum of Gaussian adiabatic density fluctuations exhibiting power and polarization on super-horizon scales, just as predicted by simple field-theoretical models of inflation. As we review below, the scale of the vacuum energy during inflation was apparently \(\sim 10^{16} \text{ GeV}\), comparable to the expected GUT scale, so CMB measurements offer us a direct window on ultra-high-energy physics.

Ever since inflation was proposed, it has been a puzzle how to integrate it with ideas in particle physics. For example, a naive GUT Higgs field would give excessive density perturbations, and no convincing concrete string-theoretical model has yet emerged. In this conceptual vacuum, models based on simple singlet scalar fields have held sway \(^1\). The simplest of these are chaotic inflation models based on exponential or power-law potentials, of which \(\phi^4\) and \(\phi^2\) are the only renormalizable examples. The WMAP collaboration has made so bold as to claim that such a \(\phi^4\) model is excluded at the 3-\(\sigma\) level \(^1\), a conclusion which would merit further support \(^3\) \(^4\). Nevertheless, it is clear that a \(\phi^2\) model would be favoured.

We reconsider in this paper the possibility that the inflaton could in fact be related to the other dramatic recent development in fundamental physics, namely the discovery of neutrino masses \(^5\). The simplest models of neutrino masses invoke heavy singlet neutrinos that give masses to the light neutrinos via the seesaw mechanism \(^6\). The heavy singlet neutrinos are usually postulated to weigh \(10^{10}\) to \(10^{15} \text{ GeV}\), embracing the range where the inflaton mass should lie, according to WMAP et al. In supersymmetric models, the heavy singlet neutrinos have scalar partners with similar masses, sneutrinos, whose properties are ideal for playing the inflaton role \(^7\). In this paper, we discuss the simplest scenario in which the lightest heavy singlet sneutrino drives inflation. This scenario constrains in interesting ways many of the 18 parameters of the minimal seesaw model for generating three non-zero light neutrino masses.

This minimal sneutrino inflationary scenario (i) yields a simple \(\frac{1}{2}m^2\phi^2\) potential with no quartic terms, with (ii) masses \(m\) lying naturally in the inflationary ballpark. The resulting (iii) spectral index \(n_s\), (iv) the running of \(n_s\) and (v) the relative tensor strength \(r\) are all compatible with the data from WMAP and other experiments \(^2\). Moreover, fixing \(m \sim 2 \times 10^{13} \text{ GeV}\) as required by the observed density perturbations (vi) is compatible with a low reheating temperature of the Universe that evades the gravitino problem \(^8\). (vii) realizes leptogenesis \(^9\) \(^10\) in a calculable and viable way, (viii) constrains neutrino model parameters, and (ix) makes testable predictions for the flavour-violating decays of charged leptons.

The main features of our scenario are the following. First, reheating of the Universe is now due to the neutrino Yukawa couplings, and therefore can be related to light neutrino masses and mixings. Secondly, the lepton asymmetry is created in direct sneutrino-inflaton decays \(^10\). There is only one parameter describing the efficiency of leptogenesis in this minimal sneutrino inflationary scenario in all leptogenesis regimes - the reheating temperature of the Universe - to

\(^1\) This argument applies a fortiori to models with \(\phi^n > 4\) potentials.
which the other relevant parameters can be related. This should be compared with the general thermal leptogenesis case [9, 11, 12, 13] which has two additional independent parameters, namely the lightest heavy neutrino mass and width. Thirdly, imposing the requirement of successful leptogenesis, we calculate branching ratios for $\mu \rightarrow e\gamma$ and $\tau \rightarrow \mu\gamma$ [14], and the CP-violating observables [15] like the electric dipole moments of the electron and muon [16]. All these leptonic observables, as well as leptogenesis, are related to the measured neutrino masses via a parametrization with a random orthogonal matrix [17]. We show that, in the minimal scenario discussed here, successful leptogenesis implies a prediction for $\mu \rightarrow e\gamma$ in a very narrow band within about one order of magnitude of the present experimental bound, whilst $\tau \rightarrow \mu\gamma$ might be somewhat further away.

Other sneutrino inflationary scenarios could be considered. For example, the inflaton might be one of the heavier singlet sneutrinos, or two or more sneutrinos might contribute to inflation, or one might play a role as a curvaton [18]. These alternatives certainly merit consideration, though they would in general be less predictive. We find it remarkable that the simplest sneutrino inflationary scenario considered here works as well as it does.

2 Chaotic Sneutrino Inflation

We start by reviewing chaotic inflation [1] with a $V = \frac{1}{2}m^2\phi^2$ potential - the form expected for a heavy singlet sneutrino - in light of WMAP [2]. Defining $M_P \equiv 1/\sqrt{8\pi G_N} \simeq 2.4 \times 10^{18}$ GeV, the conventional slow-roll inflationary parameters are

$$\epsilon \equiv \frac{1}{2}M_P^2 \left( \frac{V'}{V} \right)^2 = \frac{2M_P^2}{\phi_I^2}, \quad \eta \equiv M_P^2 \left( \frac{V''}{V} \right) = \frac{2M_P^2}{\phi_I^2}, \quad \xi \equiv M_P^4 \left( \frac{VV''}{V^2} \right) = 0,$$

where $\phi_I$ denotes the a priori unknown inflaton field value during inflation at a typical CMB scale $k$. The overall scale of the inflationary potential is normalized by the WMAP data on density fluctuations:

$$\Delta_R^2 = \frac{V}{24\pi^2M_P^2\epsilon} = 2.95 \times 10^{-9}A : A = 0.77 \pm 0.07,$$

yielding

$$V^{1/4} = M_P^4\sqrt{\epsilon} \times 24\pi^2 \times 2.27 \times 10^{-9} = 0.027M_P \times \epsilon^{1/4},$$

corresponding to

$$m_{sneutrino} = 0.038 \times M_P^3$$

in any simple chaotic $\phi^2$ inflationary model, such as the sneutrino model explore here. The number of e-foldings after the generation of the CMB density fluctuations observed by COBE is estimated to be

$$N_{COBE} = 62 - \ln \left( \frac{10^{16} \text{ GeV}}{V_{end}^{1/4}} \right) - \frac{1}{3} \ln \left( \frac{V_{end}^{1/4}}{\rho_{RH}} \right),$$

2
where $\rho_{RH}$ is the energy density of the Universe when it is reheated after inflation. The second term in (5) is negligible in our model, whereas the third term could be as large as $(-8)$ for a reheating temperature $T_{RH}$ as low as $10^6$ GeV. Conservatively, we take $N \simeq 50$. In a $\phi^2$ inflationary model, this implies

$$N = \frac{1}{4} \frac{\phi_I^2}{M^2_P} \simeq 50,$$

(6)
corresponding to

$$\phi_I^2 \simeq 200 \times M^2_P.$$

(7)

Inserting this requirement into the WMAP normalization condition (3), we find the following required mass for any quadratic inflaton:

$$m \simeq 1.8 \times 10^{13} \text{ GeV.}$$

(8)

As already mentioned, this is comfortably within the range of heavy singlet (s)neutrino masses usually considered, namely $m_N \sim 10^{10}$ to $10^{15}$ GeV.

Is this simple $\phi^2$ sneutrino model compatible with the WMAP data? The primary CMB observables are the spectral index

$$n_s = 1 - 6\epsilon + 2\eta = 1 - \frac{8 M^2_P}{\phi_I^2} \simeq 0.96,$$

(9)

the tensor-to scalar ratio

$$r \equiv \frac{A_T}{A_S} = 16\epsilon = \frac{32 M^2_P}{\phi_I^2} \simeq 0.16,$$

(10)

and the spectral-index running

$$\frac{dn_s}{d\ln k} = \frac{2}{3} \left[ (n_s - 1)^2 - 4\eta^2 \right] + 2\xi = \frac{32 M^4_P}{\phi_I^4} \simeq 8 \times 10^{-4}.$$

(11)

The value of $n_s$ extracted from WMAP data depends whether, for example, one combines them with other CMB and/or large-scale structure data. However, the $\phi^2$ sneutrino model value $n_s \simeq 0.96$ appears to be compatible with the data at the 1-$\sigma$ level. The $\phi^2$ sneutrino model value $r \simeq 0.16$ for the relative tensor strength is also compatible with the WMAP data. One of the most interesting features of the WMAP analysis is the possibility that $dn_s/d\ln k$ might differ from zero. The $\phi^2$ sneutrino model value $dn_s/d\ln k \simeq 8 \times 10^{-4}$ derived above is negligible compared with the WMAP preferred value and its uncertainties. However, $dn_s/d\ln k = 0$ appears to be compatible with the WMAP analysis at the 2-$\sigma$ level or better, so we do not regard this as a death-knell for the $\phi^2$ sneutrino model.

\footnote{In fact, we note that the favoured individual values for $n_s$, $r$ and $dn_s/d\ln k$ reported in an independent analysis\textsuperscript{4} all coincide with the $\phi^2$ sneutrino model values, within the latter’s errors!}
3 Reheating and Leptogenesis

Before addressing leptogenesis in this sneutrino model for inflation in all calculational details, we first comment on the reheating temperature $T_{RH}$ following the inflationary epoch. Assuming, as usual, that the sneutrino inflaton decays when the the Hubble expansion rate $H \sim m$, and that the expansion rate of the Universe is then dominated effectively by non-relativistic matter until $H \sim \Gamma_\phi$, where $\Gamma_\phi$ is the inflaton decay width, we estimate

$$T_{RH} = \left( \frac{90}{\pi^2 g_*} \right)^{\frac{1}{4}} \sqrt{\Gamma_\phi M_P},$$  \hspace{1cm} (12)$$

where $g_*$ is the number of effective relativistic degrees of freedom in the reheated Universe. In the minimal sneutrino inflation scenario considered here we have $\phi \equiv N_1$, $m \equiv M_{N_1}$ and

$$\Gamma_\phi \equiv \Gamma_{N_1} = \frac{1}{4\pi} (Y_\nu Y_\nu^\dagger)_{11} M_{N_1},$$  \hspace{1cm} (13)$$

where $Y_\nu$ is the neutrino Dirac Yukawa matrix. If the relevant neutrino Yukawa coupling $(Y_\nu Y_\nu^\dagger)_{11} \sim 1$, the previous choice $m = M_{N_1} \simeq 2 \times 10^{13}$ GeV would yield $T_{RH} > 10^{14}$ GeV, considerably greater than $m$ itself. Such a large value of $T_{RH}$ would be very problematic for the thermal production of gravitinos. However, it is certainly possible that $(Y_\nu Y_\nu^\dagger)_{11} \ll 1$, in which case $T_{RH}$ could be much lower, as we discuss in more detail below. Alternatively, one may consider more complicated scenarios, in which three sneutrino species may share the inflaton and/or curvaton roles between them.

We now present more details of reheating and leptogenesis. In general, inflaton decay and the reheating of the Universe are described by the following set of Boltzmann equations:

$$\frac{d\rho_\phi}{dt} = -3H \rho_\phi - \Gamma_\phi \rho_\phi,$$

$$\frac{d\rho_R}{dt} = -4H \rho_R + \Gamma_\phi \rho_\phi,$$  \hspace{1cm} (14)$$

$$H = \frac{dR}{R dt} = \sqrt{\frac{8\pi G_N (\rho_\phi + \rho_R) / 3},}$$  \hspace{1cm} (15)$$

where $\rho_\phi$ is the energy density of the inflaton field, $\rho_R$ describes the energy density of the thermalized decay products and essentially defines the temperature via

$$\rho_R = \frac{\pi^2}{30} g_* T^4,$$  \hspace{1cm} (16)$$

$H$ is the Hubble constant and $G_N$ is the Newton constant. Thus reheating can be described by two parameters, the reheating temperature $T_{RH}$, which is the highest temperature of thermal

\footnote{Even such a large value of $(Y_\nu Y_\nu^\dagger)_{11}$ would not alter significantly the $\phi^2$ sneutrino model prediction for $dn_s/d\ln k$.}
plasma immediately after reheating is completed, and the initial energy density of the inflaton field

$$\rho_\phi \simeq \frac{\pi^2 g_* T^4}{5 T_{RH}^4},$$

which determines the maximal plasma temperature in the beginning of the reheating process. In the following we use the parameter

$$z = \frac{M_{N_1}}{T}$$

to parametrize temperature.

The set of Boltzmann equations describing the inflaton decay and reheating, the creation and decays of thermal heavy neutrinos and sneutrinos, and the generation of a lepton asymmetry, is given by

$$Z \frac{d\rho_\phi}{dz} = -\frac{3\rho_\phi}{z} - \frac{\Gamma_\phi \rho_\phi}{z H},$$

$$H \frac{dY_{N_1}}{dz} = -\frac{3\Gamma_\phi \rho_\phi}{4\rho_R} \frac{Y_{N_1}}{s} - \frac{1}{s} \frac{\text{(remaining)}}{s},$$

$$H \frac{dY_{\tilde{N}_+}}{dz} = -\frac{3\Gamma_\phi \rho_\phi}{4\rho_R} \frac{Y_{\tilde{N}_+}}{s} - \frac{1}{s} \frac{\text{(remaining)}}{s},$$

$$H \frac{dY_{\tilde{N}_-}}{dz} = -\frac{3\Gamma_\phi \rho_\phi}{4\rho_R} \frac{Y_{\tilde{N}_-}}{s} - \frac{1}{s} \frac{\text{(remaining)}}{s},$$

$$H \frac{dY_{L_f}}{dz} = -\frac{3\Gamma_\phi \rho_\phi}{4\rho_R} \frac{Y_{L_f}}{s} - \frac{\Gamma_\phi \rho_\phi}{2s M_{N_1}} - \frac{1}{s} \frac{\epsilon_1}{s} \frac{\text{(remaining)}}{s},$$

$$H \frac{dY_{L_s}}{dz} = -\frac{3\Gamma_\phi \rho_\phi}{4\rho_R} \frac{Y_{L_s}}{s} - \frac{\Gamma_\phi \rho_\phi}{2s M_{N_1}} - \frac{1}{s} \frac{\epsilon_1}{s} \frac{\text{(remaining)}}{s}.$$
We are now ready to study (19-25). First we work out general results on reheating and leptogenesis in the sneutrino inflation scenario, allowing $M_{N_1}$ to vary as a free parameter. In this case, the reheating and leptogenesis efficiency is described by two parameters, namely $M_{N_1}$ and a parameter describing the decays of the sneutrino inflaton. This can be chosen to be either $\tilde{m}_1 = (Y_\nu Y_\nu^T)_{11} v^2 \sin^2 \beta / M_{N_1}$ or, more appropriately for this scenario, the reheating temperature of the Universe $T_{RH}$ given by (12). For the CP asymmetry in (s)neutrino decays, we take the maximal value for hierarchical light neutrinos, given by [20]:

$$|\epsilon^\text{max}_1(M_{N_1})| = \frac{3}{8\pi} \frac{M_{N_1} \sqrt{\Delta m^2_{\text{atm}}}}{v^2 \sin^2 \beta}.$$  \hspace{1cm} (27)

This choice allows us to study the minimal values for $M_{N_1}$ and $T_{RH}$ allowed by leptogenesis. Later, we will focus our attention on exact values of $\epsilon_1$ [21].

Solutions to (19-25) are presented in Figs. 1 and 2. We plot in Fig. 1 the parameter space in the $(M_{N_1}, \tilde{m}_1)$ plane that leads to successful leptogenesis. This parameter space has three distinctive parts with very different physics.

In the area bounded by the red dashed curve, denoted by A, leptogenesis is entirely thermal. This region has been studied in detail in [13]. Whatever lepton asymmetry is generated initially in the decay of the sneutrino inflaton is washed out by thermal effects, and the observed baryon asymmetry is generated by the out-of-equilibrium decays of thermally created singlet neutrinos.
Figure 2: The solid curve bounds the region allowed for leptogenesis in the \((T_{RH}, M_{N_1})\) plane, again obtained assuming \(Y_B > 7.8 \times 10^{-11}\) and the maximal CP asymmetry \(\epsilon_1^{\text{max}}(M_{N_1})\). In the area bounded by the red dashed curve leptogenesis is entirely thermal.

and sneutrinos. As seen in Fig. [2] in our scenario this parameter space corresponds to high \(M_{N_1}\) and high \(T_{RH}\) values.

The area B below the dashed curve and extending down to the minimum value \(M_{N_1} = 4 \times 10^6\) GeV in Fig. [1] is the region of parameter space where there is a delicate cancellation between direct lepton asymmetry production in sneutrino inflaton decays and thermal washout. This region cannot be studied without solving the Boltzmann equations numerically. However, it roughly corresponds to \(T_{RH} \sim M_{N_1}\) as seen in Fig. [2].

The area denoted by C has \(T_{RH} \ll M_{N_1}\). Since the maximal CP asymmetry scales with \(M_{N_1}\), the line presented corresponds to a constant reheating temperature. Notice that in Fig. [1] this line is terminated at \(\tilde{m}_1 = 10^{-7}\). As seen in Fig. [2] it continues linearly to high values of \(M_{N_1}\). In this area, leptogenesis is entirely given by the decays of cold sneutrino inflatons, a scenario studied previously in [10]. In this case the details of reheating are not important for our analyses. To calculate the lepton asymmetry to entropy density ratio \(Y_L = n_L/\rho\) in inflaton decays we need to know the produced entropy density

\[
s = \frac{2\pi^2}{45} g_* T_{RH}^3, \tag{28}
\]

and to take into account that inflaton dominates the Universe. In this case one obtains [10]

\[
Y_L = \frac{3}{4} \epsilon_{1} T_{RH}/M_{N_1}, \tag{29}
\]
where $\epsilon_1$ is the CP asymmetry in $\phi \equiv \tilde{N}_1$ decays. The observed baryon asymmetry of the Universe gives a lower bound on the reheating temperature $T_{RH} > 10^6$ GeV.

We consider now the most constrained scenario in which the inflaton is the lightest sneutrino, which requires $M_{N_3} > M_{N_2} > M_{N_1} \simeq 2 \times 10^{13}$ GeV. This implies that our problem is completely characterized by only one parameter, either $m_1$ or $T_{RH}$. As we see in both Figs. 1 and 2 the line for $M_{N_1} \simeq 2 \times 10^{13}$ GeV traverses both the regions A and C, the former corresponding to high $T_{RH}$, as seen in Fig. 2. However, $T_{RH}$ may also be low even in the minimal seesaw model, as seen in Fig. 2.

The cosmological gravitino problem suggests that $T_{RH} \lesssim 10^8$ GeV might be the most interesting, which would correspond to very small $\tilde{m}_1$, far away from the thermal region A and deep in the region C where leptogenesis arises from the direct decays of cold sneutrinos. We concentrate on this option here. This limit requires very small Yukawa couplings $(Y_\nu Y^\dagger_\nu)_{11} \lesssim 10^{-12}$, whilst other Yukawa couplings can be $O(1)$. This possibility may be made natural, e.g., by postulating a $Z_2$ matter parity under which only $N_1$ is odd. In this case, the relevant Yukawa couplings $(Y_\nu)_{ij}$ all vanish, but a Majorana mass for $N_1$ is still allowed. A more sophisticated model postulates a $Z_7$ discrete family symmetry with charges $Y_{FN} = (4, 0, 0)$ for the $N_i$, $(2, 1, 1)$ for the $\bar{5}$ representations of SU(5), and $(2, 1, 0)$ for the 10 representations of SU(5). Assuming a gauge-singlet field $\Phi$ with $Y_{FN} = -1$ and $\langle \Phi \rangle \equiv \epsilon$, we find $M_i = O(\epsilon, 1, 1)$ and $(Y_\nu)_{ij} = O(\epsilon^6, \epsilon^5, \epsilon^2)$, whilst the other Yukawa couplings are $O(1)$, $O(\epsilon)$ or $O(\epsilon^2)$. If $\epsilon \simeq 1/17$, the $(Y_\nu)_{ij}$ are sufficiently small for our purposes, whilst the quark and lepton mass matrices are of desirable form. Doubtless, one could construct better models with more effort, but this example serves as an existence proof for a low value of $T_{RH}$ in our scenario.

4 Leptogenesis Predictions for Lepton Flavour Violation

In this Section, we relate the results of the previous section on direct leptogenesis to light neutrino masses, and make predictions on the lepton-flavour-violating (LFV) decays. Thermal leptogenesis in this context has been extensively studied recently \[22, 23, 24\]. We first calculate neutrino Yukawa couplings using the parametrization in terms of the light and heavy neutrino masses, mixings and the orthogonal parameter matrix given in \[17\]. This allows us to calculate exactly the baryon asymmetry of the Universe, since we know the CP asymmetry $\epsilon_1$ and the reheating temperature of the Universe $T_{RH}$. For neutrino parameters yielding successful leptogenesis, we calculate the branching ratios of LFV decays.

There are 18 free parameters in the minimal seesaw model with three non-zero light neutrinos, which we treat as follows. In making Fig. 3 we have taken the values of $\theta_{12}, \theta_{23}, \Delta m^2_{12}$ and $\Delta m^2_{23}$ from neutrino oscillation experiments. We randomly generate the lightest neutrino mass in the range $0 < m_1 < 0.01$ eV and values of $\theta_{13}$ in the range $0 < \theta_{13} < 0.1$ allowed by the Chooz experiment \[25\], as we discuss later in more detail. Motivated by our previous discussion of chaotic sneutrino inflation, we fix the lightest heavy singlet sneutrino mass to be $M_1 = 2 \times 10^{13}$ GeV, and choose the following values of the heavier singlet sneutrino masses:
$M_2 = 10^{14}$ GeV or $M_2 = 5 \times 10^{14}$ GeV, and $M_3$ in the range $5 \times 10^{14}$ to $5 \times 10^{15}$ GeV, as we also discuss later in more detail. This accounts for nine of the 18 seesaw parameters.

The remaining 9 parameters are all generated randomly. These include the three light-neutrino phases - the Maki-Nakagawa-Sakata oscillation phase and the two Majorana phases. Specification of the neutrino Yukawa coupling matrix requires three more mixing angles and three more CP-violating phases that are relevant to leptogenesis, in principle. The plots in Fig. 4 are made by sampling randomly these nine parameters. We apply one constraint, namely that the generated baryon density falls within the $3 - \sigma$ range required by cosmological measurements, of which the most precise is now that by WMAP: $7.8 \times 10^{-11} < Y_B < 1.0 \times 10^{-10}$ [2].

Making predictions for LFV decays also requires some hypotheses on the parameters of the MSSM. We assume that the soft supersymmetry-breaking mass parameters $m_0$ of the squarks and sleptons are universal, and likewise the gaugino masses $m_{1/2}$, and we set the trilinear soft supersymmetry-breaking parameter $A_0 = 0$ at the GUT scale. Motivated by $g_\mu - 2$, we assume that the higgsino mixing parameter $\mu > 0$, and choose the representative value $\tan \beta = 10$. We take into account laboratory and cosmological constraints on the MSSM, including limits on the relic density of cold dark matter. WMAP provides the most stringent bound on the latter, which we assume to be dominated by the lightest neutralino $\chi_1^0$: $0.094 < \Omega_\chi h^2 < 0.129$. For $\tan \beta = 10$, the allowed domain of the $(m_{1/2}, m_0)$ plane is an almost linear strip extending from $(m_{1/2}, m_0) = (300, 70)$ GeV to $(900, 200)$ GeV [26]. For illustrative purposes, we choose $(m_{1/2}, m_0) = (800, 170)$ GeV and comment later on the variation with $m_{1/2}$.

Panel (a) of Fig. 3 presents results on the branching ratio BR for $\mu \to e\gamma$ decay. We see immediately that values of $T_{RH}$ anywhere between $2 \times 10^6$ GeV and $10^{12}$ GeV are attainable in principle. The lower bound is due to the lower bound on the CP asymmetry, while the upper bound comes from the gravitino problem. The black points in panel (a) correspond to the choice $\sin \theta_{13} = 0.0$, $M_2 = 10^{14}$ GeV, and $5 \times 10^{14}$ GeV < $M_3$ < $5 \times 10^{15}$ GeV. The red points correspond to $\sin \theta_{13} = 0.0$, $M_2 = 5 \times 10^{14}$ GeV, and $M_3 = 5 \times 10^{15}$ GeV, while the green points correspond to $\sin \theta_{13} = 0.1$, $M_2 = 10^{14}$ GeV, and $M_3 = 5 \times 10^{14}$ GeV. We see a very striking narrow, densely populated bands for BR($\mu \to e\gamma$), with some outlying points at both larger and smaller values of BR($\mu \to e\gamma$). The width of the black band is due to variation of $M_{N_3}$ showing that BR($\mu \to e\gamma$) is not very sensitive to it. However, BR($\mu \to e\gamma$) strongly depends on $M_{N_2}$ and $\sin \theta_{13}$ as seen by the red and green points, respectively. Since BR($\mu \to e\gamma$) scales approximately as $m_{1/2}^{-2}$, the lower strip for $\sin \theta_{13} = 0$ would move up close to the experimental limit if $m_{1/2} \sim 500$ GeV, and the upper strip for $\sin \theta_{13} = 0.1$ would be excluded by experiment.

Panel (b) of Fig. 3 presents the corresponding results for BR($\tau \to \mu\gamma$) with the same colour code for the parameters. This figure shows that BR($\tau \to \mu\gamma$) depends strongly on $M_{N_1}$, while the dependence on $\sin \theta_{13}$ and on $M_{N_2}$ is negligible. The numerical values of BR($\tau \to \mu\gamma$) are somewhat below the present experimental upper limit BR($\tau \to \mu\gamma$) $\sim 10^{-7}$, but we note that the results would all be increased by an order of magnitude if $m_{1/2} \sim 500$ GeV. In this case, panel (a) of Fig. 3 tells us that the experimental bound on BR($\mu \to e\gamma$) would enforce $\sin \theta_{13} \ll 0.1$, but this would still be compatible with BR($\tau \to \mu\gamma$) $> 10^{-8}$.

As a result, Fig. 3 strongly suggests that fixing the observed baryon asymmetry of the Universe
Figure 3: Calculations of BR($\mu \to e\gamma$) and BR($\tau \to \mu\gamma$) on left and right panels, respectively. Black points correspond to $\sin \theta_{13} = 0$, $M_2 = 10^{14}$ GeV, and $5 \times 10^{14}$ GeV $< M_3 < 5 \times 10^{15}$ GeV. Red points correspond to $\sin \theta_{13} = 0$, $M_2 = 5 \times 10^{14}$ GeV, and $M_3 = 5 \times 10^{15}$ GeV, while green points correspond to $\sin \theta_{13} = 0.1$, $M_2 = 10^{14}$ GeV, and $M_3 = 5 \times 10^{14}$ GeV.

for the direct sneutrino leptogenesis ($T_{RH} < 2 \times 10^{12}$ GeV $< M_{N_1}$) implies a prediction for the LFV decays provided $M_{N_2}$ and/or $M_{N_3}$ are also fixed. This observation can be understood in the case of hierarchical light and heavy neutrino masses. Consider first $\mu \to e\gamma$ for $\sin \theta_{13} = 0$. It turns out that the $N_2$ couplings dominate in $(Y_\nu Y_\nu^\dagger)_{21}$ which determines BR($\mu \to e\gamma$). Also, the $M_{N_2}$ term dominates in $\epsilon_1$ which implies $Y_B \sim (Y_\nu Y_\nu^\dagger)_{21}/\sqrt{(Y_\nu Y_\nu^\dagger)_{11}}$, because cancellations among the phases are unnatural. In the parametrization with the orthogonal matrix $R$, this implies $Y_B \sim R_{23}/R_{22}$. If fine tunings are not allowed, the requirement $T_{RH} < M_{N_1}$ fixes $R_{23}/R_{22}$ and therefore relates $Y_B$ to $\mu \to e\gamma$. For more general cases, the behaviour of BR($\mu \to e\gamma$) is more complicated and additional contributions occur. However, those new contributions tend to enhance BR($\mu \to e\gamma$), as exemplified in Fig. 3 by the green dots.

The behaviour of BR($\tau \to \mu\gamma$) is simpler. To leading order in the largest parameters, $\tau \to \mu\gamma$ depends on the $N_3$ couplings and mass, leading to $(Y_\nu Y_\nu^\dagger)_{32} \sim (Y_{\nu_3})_{33}^2 U_{33} U_{23}^\dagger$, independently of leptogenesis results.

We have to stress here that such definite predictions for LFV processes can always be avoided by fine tuning the neutrino parameters, as seen by several scattered points in Fig. 3. Points with small BR($\mu \to e\gamma$) can be systematically generated using the parametrization of $Y_\nu$ by a Hermitian matrix [27], and the predictions for the LFV decays thereby washed away. However, in this case, the $M_{N_i}$ are outputs of the parametrization, and cannot be fixed as required by the present analyses of sneutrino inflation. Therefore the parametrization [27] is not appropriate
for our leptogenesis scenario. Finally, we comment that such fine tunings are impossible in simple models of neutrino masses [24].

Another possibility for avoiding the LFV predictions is to allow the heavy neutrinos to be partially degenerate in mass, which enhances the CP asymmetries [28]. In supersymmetric models, this possibility was considered in [29].

In addition to the quantities shown in Fig. 3, we have also examined BR(τ → eγ), which is always far below the present experimental bound BR(τ → eγ) ∼ 10⁻⁷, and the electron and muon electric dipole moments. We find that de < 10⁻³³ e cm, in general, putting it beyond the foreseeable experimental reach, and |dµ/dₑ| ∼ mµ/me, rendering dµ also unobservably small.

5 Alternative Scenarios and Conclusions

We have considered in this paper the simplest sneutrino inflation scenario, in which the inflaton φ is identified with the lightest sneutrino, and its decays are directly responsible for leptogenesis. We find it remarkable that this simple scenario is not already ruled out, and have noted the strong constraints it must satisfy enable it to make strong predictions, both for CMB observables and LFV decays. These might soon be found or invalidated. In the latter case the motivation to study more complicated sneutrino inflation scenarios would be increased.

• One possibility is that inflation might have been driven by a different sneutrino, not the lightest one. In this case, the lightest sneutrino could in principle be considerably lighter than the 2 × 10¹³ GeV required for the inflaton. This would seem to make more plausible a low reheating temperature, as suggested by the gravitino problem. However, this problem is not necessarily a critical issue, as it can already be avoided in the simplest sneutrino inflation scenario, as we have seen. On the other hand, if the lightest sneutrino is not the inflaton, leptogenesis decouples from inflationary reheating, and predictivity is diminished.

• A second possibility is that two or more sneutrinos contribute to inflation. In this case, the model predictions for the CMB observables and the sneutrino mass would in general be changed.

• A related third possibility is that one or more sneutrinos might function as a curvaton, which would also weaken the CMB and sneutrino mass predictions.

For the moment, we do not see the need to adopt any of these more complicated scenarios, but they certainly merit investigation, even ahead of the probable demise of the simplest sneutrino inflation scenario investigated here.
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