Mixed inflaton and curvaton scenario with sneutrinos

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Abstract. A variation of sneutrino inflation based on χ^2 potential is considered where the inflaton and the late-decaying field are sneutrinos of different generations. The lighter, late-decaying sneutrino dilutes the gravitinos over-produced after inflaton decay and generates the matter asymmetry. It can also significantly contribute to the curvature perturbation, realizing the mixed inflaton-curvaton case. The cosmic microwave background (CMB) observables can distinguish this case from inflation with χ^2 potential, provided that the initial value of the late-decaying sneutrino is either an order of magnitude smaller or larger than the reduced Planck scale.

Keywords: cosmology of theories beyond the SM, physics of the early universe, inflation, leptogenesis

Contents

1 Introduction and review

Inflation [1] can be realized most simply by means of a scalar field χ, called the inflaton. In supersymmetric models, a well motivated inflaton candidate is one of the sneutrinos, the scalar partners of the right-handed (RH) neutrinos which determine the mass scale of the light neutrinos via the see-saw mechanism [2]. A minimal scenario where inflation is driven by a sneutrino with χ^2 potential works remarkably well [3, 4].

Among the attractive features of sneutrino inflation with χ^2 potential is that the sneutrino mass fixed from the amplitude of the primordial curvature perturbation has the same order of magnitude with the see-saw scale inferred from the light neutrino mass differences. Furthermore, this sneutrino mass ($\sim 10^{13}$ GeV) is also compatible with baryogenesis via leptogenesis [5], with the lepton asymmetry originating from the decays of the inflaton-sneutrino [6]. The predictions for the spectral index n_s and tensor to scalar ratio r are compatible with experiments including the WMAP results [7].

Sneutrino inflation predicts a reheat temperature T_r of order 10^{14} GeV if neutrino Yukawa couplings are of order unity (see section 1.1). On the other hand, to avoid the overproduction of gravitinos in the early universe, typically $T_r \lesssim 10^9$ GeV is required [8]. Avoiding the gravitino problem thus requires very small Yukawa couplings, and as a consequence the RH neutrino superpartner of the inflaton decouples from the see-saw mechanism in the sneutrino inflation and leptogenesis scenario discussed in refs. [3, 4].

It is useful to consider variations of this scenario to interpret upcoming experimental results, particularly those expected from the Planck mission [9]. Here we do not require the RH neutrino superpartner of the inflaton to have such small Yukawa couplings so that it decouples from the see-saw mechanism. Instead, we assume that there is another, lighter generation sneutrino with small Yukawa couplings, whose late decay can dilute the gravitinos produced earlier. This lighter, late-decaying sneutrino can act as a curvaton [10, 11], or can partially contribute to the primordial curvature perturbation (the mixed inflaton-curvaton case [12–17]) and also cause a second epoch of inflation [18, 19] depending on its initial value. We show that sufficient matter asymmetry can be generated while satisfying the gravitino constraint in the mixed inflaton-curvaton case and work out the predictions for the CMB observables.[1]

Inflation models are typically analyzed in terms of the slow-roll parameters, defined as:

$$\epsilon = \frac{m_P^2}{2}\left(\frac{V_\chi}{V}\right)^2 , \quad \eta = m_P^2 \frac{V_{\chi\chi}}{V} , \quad \xi^2 = m_P^4 \frac{V_\chi V_{\chi\chi\chi}}{V^2} . \tag{1.1}$$

Here $m_P \approx 2.4\times 10^{18}$ GeV is the reduced Planck scale, and the subscript 'χ' denotes derivative with respect to the inflaton χ.

The number of e-folds during inflation is given by

$$N_* \equiv \ln \frac{a(t_{end})}{a(t_*)} = \int_{t_*}^{t_{end}} H \mathrm{d}t \approx \frac{1}{m_P^2} \int_{\chi_{end}}^{\chi_*} \frac{V \mathrm{d}\chi}{V_\chi} . \tag{1.2}$$

The subscript '$_*$' implies that the values correspond to the time when the comoving wavenumber $k_* = a_* H_*$, where $(aH)^{-1}$ is the comoving Hubble length. The subscript '*end*' denotes the end of inflation, when

$$\epsilon_H \equiv 3 \frac{\dot\chi^2/2}{V + \dot\chi^2/2} = 1 . \tag{1.3}$$

The primordial curvature perturbation ζ has a nearly scale invariant spectrum. Assuming the only contribution to ζ is due to the fluctuations in the inflaton field, the amplitude of the perturbation is given by

$$\mathcal{P}_\zeta \approx \frac{V}{24\pi^2 m_P^4 \epsilon} . \tag{1.4}$$

The WMAP best fit value for the comoving wavenumber $k_* = 0.002\ \mathrm{Mpc}^{-1}$ is $\mathcal{P}_\zeta \approx 2.4\times 10^{-9}$ [7]. The spectral index n_s, the tensor to scalar ratio r and the running of the spectral index $\alpha \equiv \mathrm{d}n_s/\mathrm{d}\ln k$ are given by

$$n_s - 1 \equiv \left.\frac{\mathrm{d}\ln \mathcal{P}_\zeta}{\mathrm{d}\ln k}\right|_{k=aH} \approx -6\epsilon + 2\eta , \quad r \approx 16\epsilon , \quad \alpha \approx 16\epsilon\eta - 24\epsilon^2 - 2\xi^2 . \tag{1.5}$$

Error estimates for these expressions, which are at the leading order in slow-roll parameters, are discussed e.g. in ref. [26].

1.1 Sneutrino inflation and sneutrino dominated leptogenesis

This section is a brief review of the sneutrino inflation and leptogenesis scenario discussed in refs. [3, 4], where the minimal supersymmetric standard model is supplemented by the superpotential

$$W = \frac{1}{2} M_a N_a N_a + h_{a\alpha} N_a L_\alpha H_u . \tag{1.6}$$

Here N_a ($a = 1, 2, 3$ with the ordering $M_1 < M_2 < M_3$), L_α and H_u denote the superfields of the RH neutrinos, lepton doublets and the up-type Higgs doublet, respectively. The inflaton χ is identified with one of the three generation of sneutrinos: $\chi \equiv \sqrt{2}|\widetilde{N}_i|$. In refs. [3, 4] it is assumed that $i = 1$, while we will be considering the case $i \neq 1$ starting from section 2.

[1] For earlier work related to the curvaton mechanism see refs. [12, 20]. For a review of curvaton scenarios see ref. [21]. For discussions of sneutrino as curvaton see refs. [22–24]. For a discussion of double sneutrino inflation see ref. [25].

The inflaton mass $m_\chi \equiv M_i$ is fixed from $\mathcal{P}_\zeta$: For the potential $V = (1/2)m_\chi^2\chi^2$, eq. (1.2) gives $N_* \approx (\chi_*^2 - \chi_{end}^2)/(4m_P^2)$. Numerical calculation using the scalar field equation

$$\ddot{\chi} + 3H\dot{\chi} + m_\chi^2\chi = 0\,, \tag{1.7}$$

where $H^2 = \rho/(3m_P^2)$ and $\rho = V + \dot{\chi}^2/2$ yields $\chi_{end} \approx 1.0m_P$. We define

$$N_+ \equiv \frac{\chi_*^2}{4m_P^2} \approx N_* + \frac{\chi_{end}^2}{4m_P^2} \approx N_* + \frac{1}{4}\,. \tag{1.8}$$

Using eq. (1.4) with $\epsilon = 2m_P^2/\chi_*^2 = 1/(2N_+)$,

$$m_\chi \approx \sqrt{6\pi^2\mathcal{P}_\zeta}\cdot\frac{m_P}{N_+}\,, \tag{1.9}$$

yielding $m_\chi \approx 1.6\times 10^{13}$ GeV for $N_+ \approx 55$. From eqs. (1.1) and (1.8),

$$n_s - 1 \approx -\frac{8m_P^2}{\chi_*^2} \approx -\frac{2}{N_+}\,,\quad r \approx \frac{32m_P^2}{\chi_*^2} \approx \frac{8}{N_+}\,,\quad \alpha \approx -\frac{32m_P^4}{\chi_*^4} \approx \frac{n_s-1}{N_+}\,. \tag{1.10}$$

The inflaton-sneutrino decay width is given by

$$\Gamma_\chi = \frac{(hh^\dagger)_{ii}}{4\pi}m_\chi\,, \tag{1.11}$$

where h is the neutrino Yukawa couplings matrix. Defining the reheat temperature T_r as the temperature when $H = H_{reh} \equiv \Gamma_\chi/2$, it is given by

$$T_r \approx \left(\frac{45}{\pi^2 g_*}\right)^{1/4}\sqrt{\Gamma_\chi m_P} \approx 0.3\sqrt{\Gamma_\chi m_P}\,, \tag{1.12}$$

where we have taken the relativistic degrees of freedom $g_* \approx 200$.

The gravitino constraint on T_r depends strongly on the gravitino mass if it is unstable, ranging from 10^6–10^9 GeV for $m_{3/2} \sim 1$–10 TeV.[2] Hereafter we take the gravitino constraint to be $T_r \lesssim 10^9$ GeV, which implies $(hh^\dagger)_{ii} \lesssim 10^{-11}$. As a consequence of the suppressed Yukawa couplings, the RH neutrino superpartner of the inflaton decouples from the see-saw mechanism.[3] Since only the decoupled RH neutrino contributes to the mass of the lightest left-handed neutrino, this mass is predicted to be extremely small:

$$m_{\nu_1} \leq \widetilde{m}_i \equiv \frac{(hh^\dagger)_{ii}\langle H_u^0\rangle^2}{m_\chi} \approx \left(\frac{T_r}{m_\chi}\right)^2 2\times 10^{-3}\text{ eV} \lesssim 10^{-11}\text{ eV}\,. \tag{1.13}$$

The inflaton-sneutrino decays lead to a lepton asymmetry, given by [6]

$$Y_L \equiv \frac{n_L}{s} \approx \frac{3T_r}{4m_\chi}\epsilon \lesssim 1.5\times 10^{-10}\frac{T_r}{10^6\text{ GeV}}\,, \tag{1.14}$$

[2]For the case of stable gravitinos the constraints also depend on the next-to-lightest supersymmetric particle. For a recent analysis see ref. [27].

[3]The decoupling of a RH neutrino can help to reconcile large mixing angles with hierarchical light neutrino masses [28]. The patterns of lepton flavor violating decays associated with the decoupling assumption have been discussed in ref. [29].

where ϵ is the CP asymmetry in sneutrino decay and we have used [30]

$$\epsilon \lesssim \frac{3}{8\pi}\frac{m_\chi m_{atm}}{\langle H_u^0\rangle^2}\,. \tag{1.15}$$

The final baryon asymmetry per entropy density due to the sphaleron processes at equilibrium above the electroweak scale is given by $Y_B \approx Y_L/3$ [31]. The observed baryon asymmetry of the Universe (BAU) corresponds to $Y_B = (8.8 \pm 0.2) \times 10^{-11}$ [7]. Thus, in this example of non-thermal leptogenesis, the BAU can be obtained with $T_r \gtrsim 2 \times 10^6$ GeV.

The only free parameter relevant to the CMB observables is N_* which depends on the reheat temperature T_r logarithmically, see section 3.1. Taking leptogenesis and gravitino constraints into account, T_r is determined to within three orders of magnitude, which corresponds to an uncertainty in N_* of just about two e-folds. As a result, the predictions for the CMB observables are quite precise. Using eq. (1.10) and calculating N_* as discussed in section 3.1, we obtain $n_s = 0.963(0.964)$, $r = 0.148(0.143)$ and $\alpha = -7 \times 10^{-4}(-6 \times 10^{-4})$ for $T_r = 2 \times 10^6(10^9)$ GeV.

2 Outline of the inflaton and curvaton scenario with sneutrinos

In the sneutrino inflation and leptogenesis scenario summarized in section 1.1, it was assumed that the inflaton χ has suppressed Yukawa couplings to satisfy the gravitino constraint. We now consider an alternative situation involving the inflaton χ and a late-decaying field ϕ, which are identified with different generation RH neutrino superpartners $\widetilde{N}_i$ and $\widetilde{N}_j$, respectively. With the superpotential given by eq. (1.6), both fields χ and ϕ have quadratic potentials, and it is assumed that $m_\phi \ll m_\chi$.[4] The lighter, late-decaying ϕ field can eventually dominate the energy density of the Universe so that its decay dilutes the gravitinos produced earlier and generates the matter asymmetry [6].

We sketch possible thermal histories from the end of inflation until the decay of the ϕ field in figure 1. Using eq. (1.3) and $\chi_{end} \approx m_P$, the end of inflation corresponds to $H = H_{end} \equiv m_\chi \chi_{end}/(2m_P) \approx m_\chi/2$. The χ field then starts oscillating corresponding to a matter dominated equation of state for χ^2 potential. The χ field decays when $H \sim H_{reh} \equiv \Gamma_\chi/2$, and the ϕ field starts oscillating when $H \sim H_{osc} \equiv m_\phi/2$. We refer to the cases $m_\phi < \Gamma_\chi$ and $m_\phi > \Gamma_\chi$ as case 1 and case 2 respectively. The energy density of the ϕ field ρ_ϕ equals half the total energy density when $H \equiv H_e$ and the ϕ field dominates afterwards. Finally, the ϕ field decays when $H \sim H_d \equiv \Gamma_\phi/2$.

Different inflationary scenarios can occur depending on ϕ_*, the initial field value when the comoving wavenumber k_* exits horizon during inflation. As discussed in section 3, for $\phi_* \ll 0.1 m_P$ the ϕ field is the curvaton while $\phi_* \sim 0.1 m_P$ corresponds to the mixed inflaton-curvaton case, that is, both $\delta\chi$ and $\delta\phi$ significantly contribute to ζ. For $\phi_* \gtrsim m_P$ (case b), the ϕ field dominates before it starts oscillating and leads to a second epoch of inflation, which can start either after reheating (case 1b) or during χ oscillations (case 2b). In either case, $\phi_* \sim 10 m_P$ again corresponds to the mixed inflaton-curvaton case. For convenience we will refer to the late-decaying ϕ field as the curvaton even though it only partially contributes to ζ in general.

[4] The scalar potential is generally altered by supergravity corrections, for a review see ref. [26]. It is nevertheless also possible for scalar fields to have quadratic potentials up to field values much greater than m_P in supergravity, with specific non-minimal Kähler potentials [32] or with a shift-symmetric Kähler potential [33].

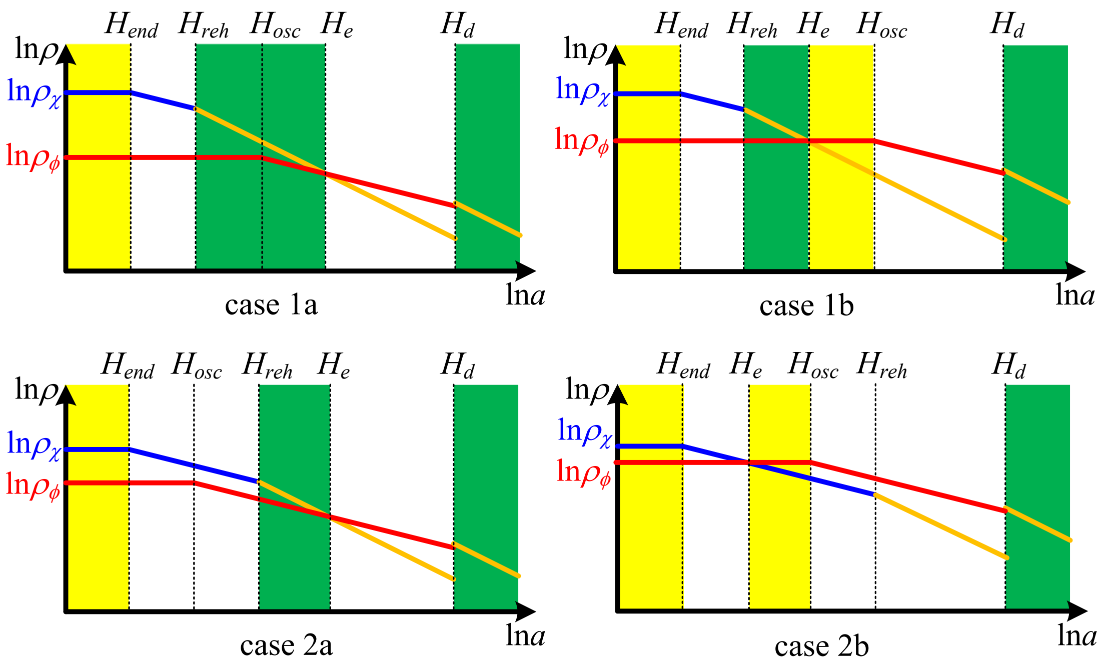

Figure 1. Possible thermal histories. Inflationary epochs and radiation dominated epochs are shaded yellow (light) and green (dark), respectively. In the unshaded regions oscillations of either field dominate, corresponding to matter dominated equation of state. The orange (light) segments indicate the energy density in radiation. Case 1 and case 2 correspond to $m_\phi < \Gamma_\chi$ and $m_\phi > \Gamma_\chi$ respectively, with a: $\phi_* \lesssim m_P$, b: $\phi_* \gtrsim m_P$.

We can estimate H_e by setting $\rho_\phi = \rho_\chi$ for case 2b, and $\rho_\phi = \rho_r$ in the other cases, ρ_r being the energy density in radiation. The ϕ field stays almost constant at ϕ_* until $H = H_{osc}$ for $\phi_* \lesssim m_P$ (case a). For $\phi_* \gtrsim m_P$ (case b) it stays almost constant until $H = H_e$, and decreases to $\phi \approx m_P$ by the time the second epoch of inflation ends at $H \approx H_{osc}$. Neglecting the change in ρ_ϕ until H_{osc} and H_e for case a and case b respectively, we obtain

$$\text{for case a}: \ H_e \sim \left(\frac{\phi_*}{m_P}\right)^4 \min(m_\phi, \Gamma_\chi)\,, \quad \text{for case b}: \ H_e \sim \left(\frac{\phi_*}{m_P}\right) m_\phi\,. \tag{2.1}$$

2.1 Constraints on the curvaton and inflaton parameters

The curvaton initial value: A lower bound on ϕ_* follows from the condition that ϕ dominates before it decays, so that the pre-existing gravitinos can be diluted. From setting $H_e = H_d$ we obtain

$$\text{for case a}: \ \left(\frac{\phi_*}{m_P}\right)^4 \gtrsim \frac{\Gamma_\phi}{\min(m_\phi, \Gamma_\chi)}\,, \quad \text{for case b}: \ \frac{\phi_*}{m_P} \gtrsim \frac{\Gamma_\phi}{m_\phi}\,. \tag{2.2}$$

However, for sufficient dilution of gravitinos ϕ must dominate much earlier. Taking the gravitino constraint on the reheat temperature to be $T_r \lesssim 10^9$ GeV and recalling that the thermal abundance of gravitinos are proportional to T_r, the required dilution factor $\Delta_{\rm req} \sim T_r/(10^9\ \text{GeV})$. The dilution factor due to ϕ decay is $\Delta \approx 1.8 g_*^{1/4} m_\phi (n_\phi/s)/\sqrt{\Gamma_\phi m_P}$ [34].

Estimating this dilution factor for each thermal history sketched in figure 1, lower bounds on ϕ_* corresponding to $\Delta > \Delta_{\rm req}$ are obtained as follows:

$$\begin{aligned}
&\text{for case a}: && \left(\frac{\phi_*}{m_P}\right)^4 \gtrsim \frac{\Gamma_\chi}{\min(m_\phi, \Gamma_\chi)} \cdot \frac{\Gamma_\phi}{1\ \text{GeV}} \geq \frac{\Gamma_\phi}{1\ \text{GeV}}\,, \\
&\text{for case 1b}: && \left(\frac{m_P}{\phi_*}\right)^3 \exp\left[\frac{3}{2}\left(\left(\frac{\phi_*}{m_P}\right)^2 - 1\right)\right] \gtrsim \frac{\Gamma_\chi}{m_\phi} \cdot \frac{\Gamma_\phi}{1\ \text{GeV}}\,, \\
&\text{for case 2b}: && \left(\frac{\phi_*}{m_P}\right)^4 \exp\left[\frac{9}{8}\left(\left(\frac{\phi_*}{m_P}\right)^2 - 1\right)\right] \gtrsim \frac{\Gamma_\phi}{1\ \text{GeV}}\,.
\end{aligned} \tag{2.3}$$

The curvaton decay width and mass: Assuming the curvaton dominates before it decays, the reheat temperature after the decay is obtained as in eq. (1.12): $T_d \approx 0.3\sqrt{\Gamma_\phi m_P}$. The lepton asymmetry created by the decays of ϕ is given by eq. (1.14) with T_r and m_χ replaced by T_d and m_ϕ. The constraints from sneutrino dominated leptogenesis and gravitino over-production (2×10^6 GeV $\lesssim T_d \lesssim 10^9$ GeV) correspond to 10^{-5} GeV $\lesssim \Gamma_\phi \lesssim 10$ GeV.

Throughout the paper we assume that the curvaton mass $m_\phi \ll m_\chi$. As for a lower bound, generating the BAU via sneutrino dominated leptogenesis requires $m_\phi \gtrsim T_d \gtrsim 2\times 10^6$ GeV. To avoid thermalization of the condensate, $\widetilde{m}_j \lesssim 2\times 10^{-3}$ eV is required [22, 23]. Since

$$m_{\nu_1} \leq \widetilde{m}_j \equiv \frac{(hh^\dagger)_{jj}\langle H_u^0\rangle^2}{m_\phi} \approx \left(\frac{T_d}{m_\phi}\right)^2\ 2\times 10^{-3}\ \text{eV}\,, \tag{2.4}$$

we see that this condition is also satisfied for $m_\phi \gtrsim T_d$. Note that part of the asymmetry can be washed out if $m_\phi \sim T_d$, but for $\widetilde{m}_j \lesssim 2\times 10^{-3}$ eV the washout is weak and enough asymmetry can remain [35].

Another constraint on Γ_ϕ and m_ϕ comes from the asymmetry created by a variation of the Affleck-Dine (AD) mechanism [36], discussed in ref. [37]. The asymmetry is produced due to the rotation of the condensate as in the AD mechanism, but the asymmetry that survives the decay of sneutrinos depends on supersymmetry breaking as in soft leptogenesis [38]. Ref. [23] estimates the asymmetry resulting from the soft supersymmetry breaking B-term for the sneutrino and concludes that $m_\phi \gtrsim 10^8$ GeV and $m_{\nu_1} \lesssim 10^{-8}$ eV is required for sneutrino dominated leptogenesis to work. On the other hand, it is shown in ref. [37] that the asymmetry from the B-term is negligible by itself, and including the more important thermal effects results in an asymmetry that for $m_\phi \gtrsim T_d$ is given by

$$\frac{n_L}{s} \sim 10^{-8}\left(\frac{\Gamma_\phi}{10^{-5}\ \text{GeV}}\right)\left(\frac{1\ \text{TeV}}{B}\right)\left(\frac{T_d}{m_\phi}\right)\left|\exp\left(-\frac{m_\phi}{2T_d}\right) - \frac{0.09T_d^2}{m_\phi^2}\right|\,. \tag{2.5}$$

It is thus possible for the BAU to be generated by this mechanism with $m_\phi \gtrsim T_d \gtrsim 10^{5.5}$ GeV. However, baryon isocurvature perturbations arise if the ϕ field starts oscillating before it dominates (case a) [39, 23]. For these isocurvature perturbations to remain below the observational bounds, either the asymmetry given by eq. (2.5) should be subdominant (which requires $m_\phi \gtrsim 4T_d$) or the contribution from the ϕ field to the curvature perturbation should be subdominant. As discussed in section 3, keeping this contribution below 10% requires $\phi_* \gtrsim 0.3 m_P$.

The inflaton decay width and mass: The inflaton decay width Γ_χ is given by eq. (1.11). Using eq. (1.13), it can be expressed as follows:

$$\Gamma_\chi = \frac{\widetilde{m}_i m_\chi^2}{4\pi \langle H_u^0 \rangle^2} \sim \left(\frac{\widetilde{m}_i}{0.05 \text{ eV}} \right) \left(\frac{m_\chi}{10^{13} \text{ GeV}} \right)^2 10^{10} \text{ GeV} . \tag{2.6}$$

If the contribution of $\delta\phi$ to ζ is negligible, $m_\chi \approx 1.6 \times 10^{13}$ GeV as mentioned in section 1.1. As will be explained in section 3.1, m_χ depends on ϕ_* in the mixed inflaton-curvaton case. Applying the WMAP and gravitino constraints on ϕ_* and using eq. (3.6), m_χ is found to vary in the range 0.9–2.2$\times 10^{13}$ GeV, see figure 3.

3 Inflationary predictions

3.1 Calculating the power spectrum and the number of e-folds

Using the δN formalism [18, 40, 16], the primordial curvature perturbation ζ can be written as

$$\zeta = \delta N_{tot} \approx \frac{\partial N_{tot}}{\partial \chi} \delta\chi_* + \frac{\partial N_{tot}}{\partial \phi} \delta\phi_* , \tag{3.1}$$

where N_{tot} is the number of e-folds from horizon exit (of the scale corresponding to the comoving wavenumber k_*) to some final time t_f well after the curvaton has decayed, when $H = H_f \ll \Gamma_\phi$. We can separate N_{tot} into two parts $N_{tot} = N_* + N$, where N_* is the number of e-folds from horizon exit to the end of inflation and N is the number of e-folds from the end of inflation to t_f. Since N does not depend on χ,

$$\frac{\partial N_{tot}}{\partial \chi} = \frac{\partial N_*}{\partial \chi} \approx \frac{V}{m_P^2 V_\chi} , \tag{3.2}$$

where in the last step we have used eq. (1.2). Similarly, since N_* does not depend on ϕ, $\partial N_{tot}/\partial\phi = \partial N/\partial\phi \equiv N_\phi$. Thus,

$$\zeta \approx \frac{V}{m_P^2 V_\chi} \delta\chi_* + N_\phi \delta\phi_* . \tag{3.3}$$

Assuming $\delta\chi_*$ and $\delta\phi_*$ to be uncorrelated, the power spectrum of the perturbation is then

$$\mathcal{P}_\zeta \approx \left(\frac{V^2}{m_P^4 V_\chi^2} + N_\phi^2 \right) \left(\frac{H}{2\pi} \right)^2 . \tag{3.4}$$

Defining $y \equiv 2 m_P^2 N_\phi^2 \epsilon$, this equation can be written as

$$\mathcal{P}_\zeta \approx \frac{(1+y)V}{24\pi^2 m_P^4 \epsilon} . \tag{3.5}$$

The contribution of the ϕ field to the curvature perturbation is negligible for $y \ll 1$ whereas $y \gg 1$ corresponds to the curvaton limit. For $y \sim 1$ the mixed inflaton-curvaton case is realized. From eqs. (3.5) and (1.8) the inflaton mass is given as

$$m_\chi \approx \sqrt{\frac{6\pi^2 \mathcal{P}_\zeta}{1+y}} \cdot \frac{m_P}{N_+} . \tag{3.6}$$

Since y depends on N_* as well as N_ϕ, we now discuss how to calculate the e-fold numbers N_* and N. Using the definition of $N_* \equiv \ln(a_{end}/a_*)$ where $k_* = a_* H_*$ at horizon exit, we can relate N_* to the current scale factor a_0 and Hubble parameter H_0 as follows:

$$N_* = -\ln \frac{k_*}{a_0 H_0} + \ln \frac{a_{end}}{a_f} + \ln \frac{a_f}{a_0} + \ln \frac{H_*}{H_0}\,. \tag{3.7}$$

In this expression the first term at the right hand side is fixed from $k_* = 0.002$ Mpc^{-1}, the second term is $-N$ and the last term can be expressed in terms of m_χ and N_*. For the third term note that any significant entropy production after t_f would dilute the $B-L$ asymmetry created by the decays of ϕ. Therefore assuming no significant entropy production,

$$\ln \frac{a_f}{a_0} = \frac{1}{3} \ln \frac{s_0}{s_f} = \frac{1}{3} \ln \frac{g_{*s0} T_0^3}{g_{*s} T_f^3}\,. \tag{3.8}$$

Using $\rho_f = 3H_f^2 m_P^2 = (\pi^2/30) g_* T_f^4$ and taking the relativistic degrees of freedom $g_* = g_{*s} = 200$ we obtain

$$N_* - \frac{1}{2} \ln N_* \approx 64.3 + \ln \frac{m_\chi}{m_P} + \frac{1}{2} \ln \frac{m_P}{H_f} - N\,. \tag{3.9}$$

To estimate N, we can add the number of e-folds in each matter dominated, radiation dominated or inflationary epoch between the end of inflation t_{end} corresponding to $H = H_{end}$ and the final time t_f corresponding to $H = H_f$ assuming the transitions between the epochs to be sudden. For instance, in the case of sneutrino inflation with no late-decaying ϕ field, the universe has a matter dominated equation of state (for χ^2 potential) between H_{end} and H_{reh}, and radiation dominates after $H = H_{reh}$. Therefore $N \approx (2/3) \ln(H_{end}/H_{reh}) + (1/2) \ln(H_{reh}/H_f)$. Using this with eqs. (3.9) and (1.10), we obtain the values of n_s, r and α given in section 1.1.

In the presence of the late-decaying ϕ field, four possible thermal histories were discussed in section 2. For case 1a, there are alternating matter-radiation-matter-radiation dominated epochs between H_{end} and H_f (see figure 1) so that

$$N \approx \frac{2}{3} \ln \frac{H_{end}}{H_{reh}} + \frac{1}{2} \ln \frac{H_{reh}}{H_e} + \frac{2}{3} \ln \frac{H_e}{H_d} + \frac{1}{2} \ln \frac{H_d}{H_f}\,. \tag{3.10}$$

Using $H_{end} \approx m_\chi/2$, $H_{reh} \equiv \Gamma_\chi/2$, $H_e \sim (\phi_*/m_P)^4 m_\phi$ and $H_d \equiv \Gamma_\phi/2$, the result shown in table 1 is obtained. For case 1b, the ϕ field dominates before it starts oscillating, leading to a second inflationary epoch between H_e and H_{osc} lasting $\approx \phi_*^2/(4m_P^2)$ e-folds. The results are similar for case 2. Note that both for case 1 and case 2, $N_\phi \approx 2/(3\phi_*)$ in the limit $\phi_* \ll m_P$, and $N_\phi \approx \phi_*/(2m_P^2)$ in the limit $\phi_* \gg m_P$. These results were also obtained in refs. [14–16], where only case 1 was considered.

For a more accurate calculation, we numerically solve the following background equations, from t_{end} to $t_f \gg \Gamma_\phi^{-1}$:

$$\dot{\rho_\chi} + 3H\rho_\chi = -\Gamma_\chi \rho_\chi\,, \quad \ddot{\phi} + (3H + \Gamma_\phi)\dot{\phi} + m_\phi^2 \phi = 0\,, \quad \dot{\rho_r} + 4H\rho_r = \Gamma_\chi \rho_\chi + \Gamma_\phi \rho_\phi\,, \tag{3.11}$$

where $\rho_\phi = \dot{\phi}^2/2 + (1/2) m_\phi^2 \phi^2$ and $H^2 = (\rho_\chi + \rho_\phi + \rho_r)/(3m_P^2)$. The initial conditions are taken as follows:

$$\rho_\chi(t_{end}) = \frac{3}{4} m_\chi^2 m_P^2\,, \quad \rho_r(t_{end}) = 0\,, \quad \phi(t_{end}) = \phi_*\,, \quad \dot{\phi}(t_{end}) = -\frac{m_\phi^2 \phi_*}{3H_{end}}\,. \tag{3.12}$$

	Case a ($\phi_* \lesssim m_P$)	Case b ($\phi_* \gtrsim m_P$)
Case 1 ($m_\phi < \Gamma_\chi$)	$N \approx \frac{2}{3}\ln\frac{\phi_*}{m_P} + \frac{1}{6}\ln\frac{m_\chi^4 m_\phi}{\Gamma_\chi \Gamma_\phi H_f^3}$	$N \approx \frac{\phi_*^2}{4m_P^2} - \frac{1}{2}\ln\frac{\phi_*}{m_P} + \frac{1}{6}\ln\frac{m_\chi^4 m_\phi}{\Gamma_\chi \Gamma_\phi H_f^3}$
Case 2 ($m_\phi > \Gamma_\chi$)	$N \approx \frac{2}{3}\ln\frac{\phi_*}{m_P} + \frac{1}{6}\ln\frac{m_\chi^4}{\Gamma_\phi H_f^3}$	$N \approx \frac{\phi_*^2}{4m_P^2} - \frac{2}{3}\ln\frac{\phi_*}{m_P} + \frac{1}{6}\ln\frac{m_\chi^4}{\Gamma_\phi H_f^3}$

Table 1. The approximate number of e-folds N from the end of inflation t_{end} to a final time t_f, for the four thermal histories discussed in section 2.

To calculate N and N_ϕ, we first estimate m_χ using eqs. (3.6), (3.9) and the approximate expressions for N given in table 1. Using eqs. (3.11) and (3.12) we evaluate N numerically for different values of ϕ_* and interpolate to obtain N as a function of ϕ_*. We then recalculate m_χ and iterate this procedure.

3.2 Observational quantities

In the sneutrino inflaton-curvaton scenario we have outlined, we have assumed that the late-decaying curvaton dominates the Universe before it decays. Non-Gaussianities are then not expected to be large [11, 16]. As discussed in section 2.1, there could be isocurvature perturbations at an observable level if both $\phi_* \lesssim 0.3 m_P$ and the matter asymmetry mostly originates from a variation of the AD mechanism. Since this is not a general prediction, our discussion of observational quantities will focus on the usual parameters n_s and r. We will also briefly comment on α.

In the presence of the late-decaying ϕ field, the following expressions are obtained using eq. (3.5) [14, 16]:

$$n_s - 1 \approx -2\epsilon + \frac{2\eta - 4\epsilon}{1+y}\,, \quad r \approx \frac{16\epsilon}{1+y}\,, \quad \alpha \approx 4\epsilon(\eta - 2\epsilon) + \frac{12\epsilon\eta - 16\epsilon^2 - 2\xi^2}{1+y}\,. \tag{3.13}$$

For inflation with χ^2 potential, $\xi^2 = 0$ and from eq. (1.8), $\epsilon = \eta = 1/(2N_+)$. In terms of $N_+ \approx N_* + 1/4$ and y we have

$$n_s - 1 \approx -\frac{1}{N_+}\left(\frac{2+y}{1+y}\right)\,, \quad r \approx \frac{8}{N_+(1+y)}\,, \quad \alpha \approx \frac{n_s - 1}{N_+}\,. \tag{3.14}$$

As discussed in section 3.1, the values of the decay widths Γ_χ, Γ_ϕ and the curvaton mass m_ϕ only have a small (logarithmic) effect on the number of e-folds. Thus, values of n_s, r and α depend mostly on the initial value ϕ_*. For a qualitative discussion of the results it is convenient to define another parameter $N_0 \equiv N_+ + \phi_*^2/(4m_P^2)$. Since the ϕ field leads to a second inflationary epoch lasting $\approx \phi_*^2/(4m_P^2)$ e-folds for $\phi_* \gtrsim m_P$, N_0 is approximately the total number of e-folds during the two inflationary epochs.

For $\phi_* \lesssim m_P$ (case a), $N_0 \approx N_+ \approx 55$. Using $N_\phi \approx 2/(3\phi_*)$ and $y = m_P^2 N_\phi^2/N_+$, we see that the curvaton limit $y \gg 1$ applies if $\phi_*/m_P \ll 2/(3\sqrt{N_+}) \approx 0.09$. In this limit

$$n_s - 1 \approx -\frac{1}{N_+}\,, \quad r \approx \frac{8}{yN_+} \approx \frac{18\phi_*^2}{m_P^2}\,, \quad \alpha \approx -\frac{1}{N_+^2}\,. \tag{3.15}$$

For $\phi_* \sim m_P$, the standard predictions of inflation with χ^2 potential – given by eq. (1.10) – are recovered since $y \ll 1$. Finally, for $\phi_* \gtrsim m_P$ (case b), $N_+ \approx N_0 - \phi_*^2/(4m_P^2)$ while N_0

remains approximately constant. Using $N_\phi \approx \phi_*/(2m_P^2)$ we obtain $y \approx \phi_*^2/(4m_P^2 N_+)$ and $yN_+ \approx N_0 - N_+$. Expressing eq. (3.14) in terms of N_0 and N_+ yields

$$n_s - 1 \approx -\frac{1}{N_+} - \frac{1}{N_0}\,, \quad r \approx \frac{8}{N_0}\,, \quad \alpha \approx \frac{n_s - 1}{N_+}\,. \tag{3.16}$$

It follows from these expressions that the spectrum is more red-tilted for larger values of ϕ_*, whereas the tensor to scalar ratio r remains essentially constant. Note that while $\delta\phi$ can partially contribute to ζ, the curvaton limit $y \gg 1$ is not possible for case b, since cosmological scales do not exit the horizon during the initial epoch of inflation driven by χ if $\phi_* \gtrsim 14 m_P$.[5]

We now consider how the results depend on the decay widths Γ_χ, Γ_ϕ and the curvaton mass m_ϕ. N_0 increases with Γ_ϕ and also if $\Gamma_\chi > m_\phi$, since the radiation dominated epochs last longer. Following the discussion of the gravitino constraint in section 2, for the high N_0 case we take $\Gamma_\phi = 10$ GeV corresponding to a reheat temperature $T_d \sim 10^9$ GeV. We also take $\Gamma_\chi = 10^{12}$ GeV and $m_\phi = 10^{11}$ GeV corresponding to case 1.[6] For the low N_0 case we take $\Gamma_\phi = 10^{-5}$ GeV, since it is difficult to obtain sufficient matter asymmetry for lower values. Minimizing N_0 requires $\Gamma_\chi < m_\phi$, to be specific we take $\Gamma_\chi = 10^9$ GeV and $m_\phi = 10^{11}$ GeV, corresponding to case 2. From table 1 we see that there is about three e-folds difference between the two cases.

Numerical results for the two cases are displayed in figure 2 and figure 3. Note that although the tensor to scalar ratio r becomes negligible in the small ϕ_* limit, the leptogenesis constraint $\Gamma_\phi \gtrsim 10^{-5}$ GeV together with the gravitino constraint eq. (2.3) requires $\phi_* \gtrsim 10^{-5/4} m_P$. This implies $r \gtrsim 0.04$, which can be observed by the Planck satellite [9, 42]. For the high N_0 case the gravitino constraint corresponds to $\phi_* \gtrsim 2.4 m_P$.

To summarize, in the presence of the late-decaying ϕ field, the predictions for n_s and r can be distinguished from the predictions of inflation with χ^2 potential for values of ϕ_* which are either $\sim 0.1 m_P$ or $\sim 10 m_P$, corresponding to the mixed inflaton-curvaton case. (The curvaton limit $\phi_* \ll 0.1 m_P$ is disfavored by the gravitino constraint.) For $\phi_* \sim 0.1 m_P$, $n_s \approx 0.97$ and $r < 0.1$ yet large enough to be observable. For $\phi_* \sim 10 m_P$, $n_s \lesssim 0.96$ and $r = 0.14$–0.15.

In the mixed inflaton-curvaton case the single field consistency relation $n_T = -r/8$ gets modified to [14]

$$n_T = -\frac{(1+y)r}{8}\,, \tag{3.17}$$

implying a more red-tilted tensor spectrum. Although testing the single-field consistency relation and observing a deviation from it is beyond the forecasted accuracy of the Planck mission, it might be possible with future CMB observations provided $r \gtrsim 0.1$ [44].

The running of the spectral index α remains small in general, although it is enhanced for $\phi_* \gg m_P$. The WMAP constraints on n_s and r are satisfied for $\phi_* \lesssim 9.6 m_P$, which implies $\alpha \gtrsim -1.6 \times 10^{-3}$. This running could perhaps be observed by future galaxy surveys and 21 cm experiments [45].

[5] The curvaton potential must be different from quadratic for the inflating curvaton scenario to be realized [41].

[6] Since $m_\phi \sim 100 T_d$, the asymmetry given by eq. (2.5) is suppressed to an order of magnitude compatible with the observed BAU.

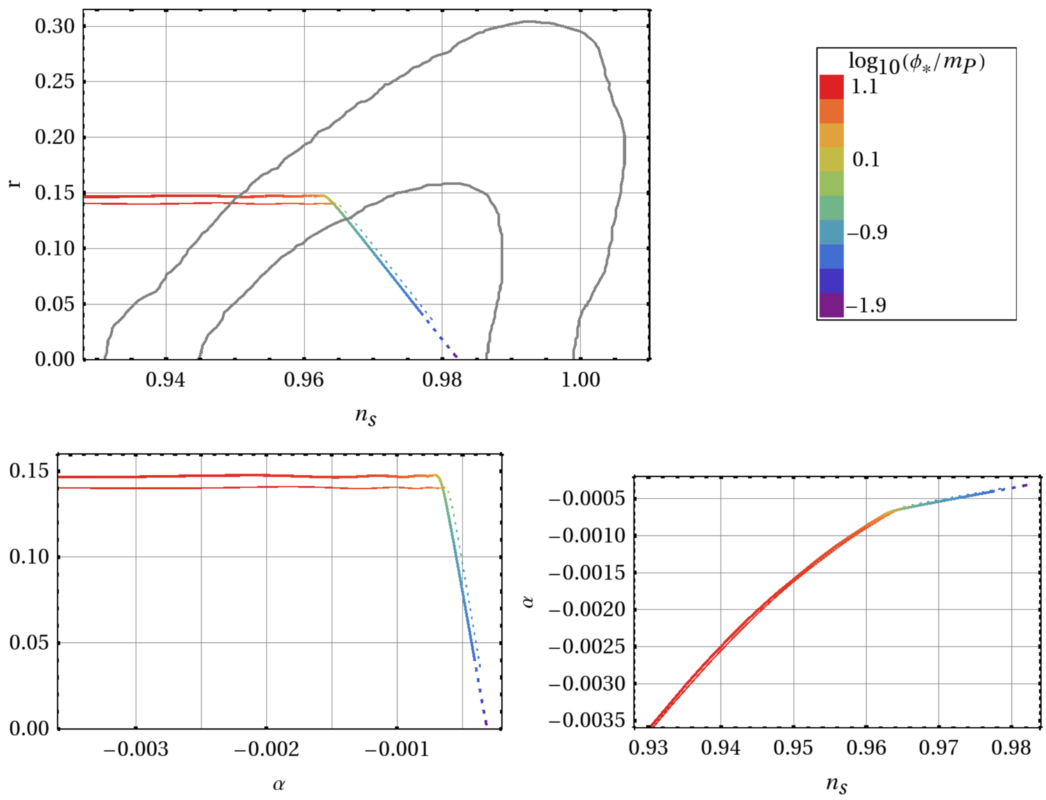

Figure 2. r vs n_s, r vs α and α vs n_s for the low N_0 case (thick curves) and the high N_0 case (thin curves). The gravitino constraint eq. (2.3) is not satisfied in the dashed segments. The r vs n_s plot also displays the WMAP contours as obtained with the WMAP Cosmological Parameter Plotter (see ref. [43], model: lcdm+sz+lens+tens, data: wmap7+bao+h0).

4 Conclusion and discussion

In the simple yet successful scenario of sneutrino inflation based on χ^2 potential, the sneutrino driving inflation can also generate the matter asymmetry via non-thermal leptogenesis [3, 4]. We have considered a variation of this scenario, by assuming that the sneutrino driving inflation is distinct from the late-decaying sneutrino. The lighter, late-decaying sneutrino can partially contribute to the curvature perturbation and alter the predictions for the CMB observables. The late decay of this sneutrino also dilutes the gravitinos produced earlier and generates the matter asymmetry.

In this mixed inflaton-curvaton scenario, since the late-decaying sneutrino is lighter its Yukawa couplings and therefore the lightest neutrino mass m_{ν_1} do not have to be as small as in the original version of sneutrino inflation, which requires $m_{\nu_1} \lesssim 10^{-11}$ eV to satisfy the gravitino constraint. However m_{ν_1} is still constrained to be $\lesssim 2 \times 10^{-3}$ eV (see section 2.1).

Considering the CMB observables the mixed inflaton-curvaton scenario is less predictive, since n_s and r depend on ϕ_*, the initial value of the late-decaying sneutrino. For $\phi_* \sim m_P$ the predictions are the same as inflation with χ^2 potential. However, the predictions change significantly for values of ϕ_* which are either $\sim 0.1 m_P$ or $\sim 10 m_P$. In the first case $n_s \approx 0.97$ and $r < 0.1$ yet large enough to be observable, whereas in the latter case $n_s \lesssim 0.96$ and $r = 0.14$–0.15.

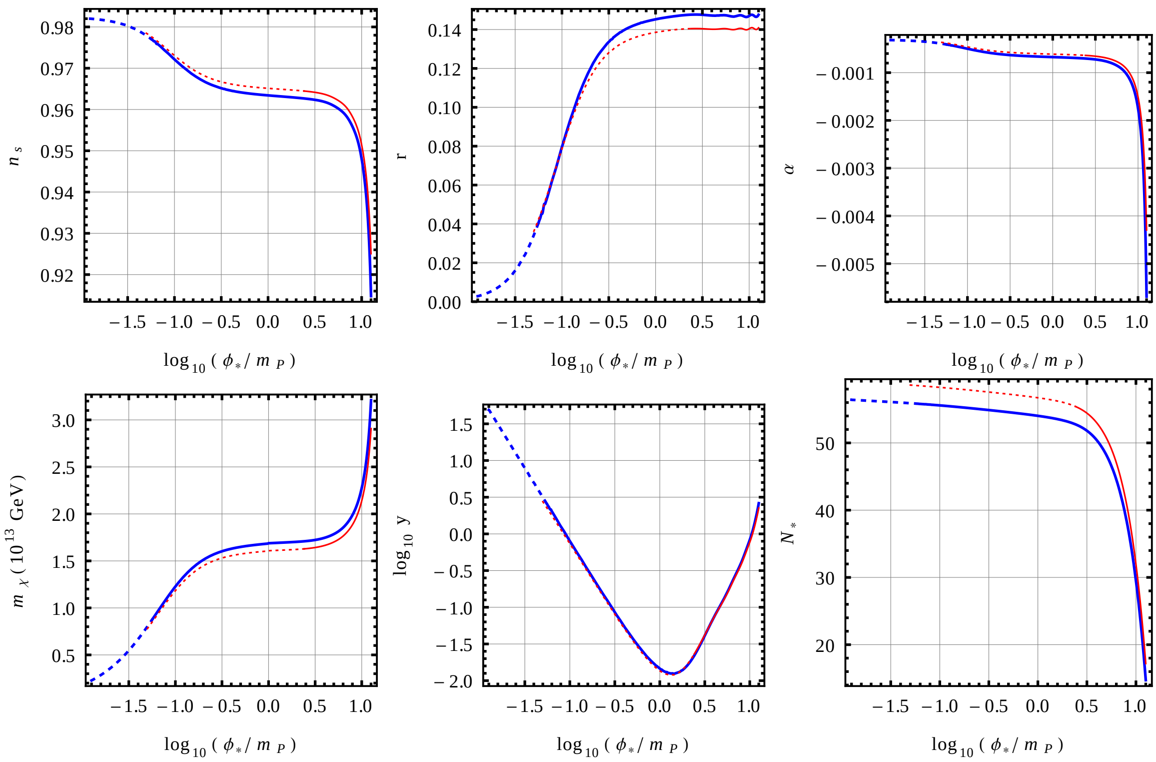

Figure 3. n_s, r, α, m_χ, y and N_* vs ϕ_* for the low N_0 case (thick blue curves) and the high N_0 case (thin red curves). The gravitino constraint eq. (2.3) is not satisfied in the dashed segments.

The original version of sneutrino inflation fixes the mass of one RH neutrino but provides no constraints on the other two, except that they should not decouple from the see-saw mechanism. Whereas in the mixed inflaton-curvaton scenario the inflaton sneutrino is distinct from the late-decaying sneutrino, so it is possible to determine the mass of the former using the CMB observables and put a lower limit on the mass of the latter from the observed matter asymmetry. If future CMB data remain consistent with this scenario, it would be of interest to embed it in a more predictive model connecting CMB observables to low-energy leptonic observables.

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