Sneutrino Hybrid Inflation

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Abstract. We review the scenario of sneutrino hybrid inflation, where one of the singlet sneutrinos, the superpartners of the right-handed neutrinos, plays the role of the inflaton. In a minimal model of sneutrino hybrid inflation, the spectral index is given by $n_s \approx 1 + 2\gamma$. With $\gamma = 0.025 \pm 0.01$ constrained by WMAP, a running spectral index $|dn_s/d\ln k| \ll |\gamma|$ and a tensor-to-scalar ratio $r \ll \gamma^2$ are predicted. Small neutrino masses arise from the seesaw mechanism, with heavy masses for the singlet (s)neutrinos generated by the vacuum expectation value of the waterfall field after inflation. The baryon asymmetry of the universe can be explained by non-thermal leptogenesis via sneutrino inflaton decay, with low reheat temperature $T_{RH} \approx 10^6$ GeV.

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INTRODUCTION

The interface between early universe cosmology and particle physics provides many challenges. Open fundamental questions in this context include the identity of the scalar field responsible for inflation \[1\], in order to solve flatness and horizon problems of the early universe, and the origin of the observed baryon asymmetry.

The experimental discovery of neutrino mass and mixing, when combined with the ideas of the see-saw mechanism and supersymmetry, gives a new perspective on both of these challenges. In order to generate the observed neutrino masses within a see-saw extended version of the Minimal Supersymmetric Standard Model (MSSM), right-handed neutrinos (together with their scalar superpartners, the singlet sneutrinos) are typically introduced and small neutrino masses arise naturally from the see-saw mechanism. In see-saw scenarios, the out-of-equilibrium decay of these right-handed (s)neutrinos in the early universe can generate the observed baryon asymmetry \[2\]. Among the particles of this extended MSSM, the singlet sneutrinos become attractive candidates for playing the role of the inflaton. Motivated by such considerations, the possibility of chaotic (large field) inflation with a sneutrino inflaton decay has been proposed \[3\]. An alternative to chaotic inflation is hybrid inflation \[4\], which, in contrast to chaotic inflation, involves field values well below the Planck scale and is thereby promising for connecting inflation to particle physics.

In this talk we review the scenario proposed in \[5\], that one of the singlet sneutrinos $\tilde{N}_i$, where $i = 1,2,3$ is a family index, plays the role of the inflaton field of hybrid inflation. We present a minimal model of sneutrino hybrid inflation, investigate prospects of generating the baryon asymmetry of our universe via non-thermal leptogenesis \[2,6\] and discuss how future observations can distinguish sneutrino hybrid inflation from scenarios of chaotic sneutrino inflation.
SUPERPOTENTIAL FOR SNEUTRINO HYBRID INFLATION

In order to illustrate how sneutrino hybrid inflation can be realized, let us consider the following minimal superpotential

$$W = \kappa \hat{S} \left( \frac{\hat{\phi}^4}{M'^2} - M^2 \right) + \frac{(\lambda_N)_{ij}}{M_*} \hat{N}_i \hat{N}_j \hat{\phi} \hat{\phi} + \ldots.$$  \hspace{1cm} (1)

\(\kappa\) and \((\lambda_N)_{ij}\) are dimensionless Yukawa couplings and \(M, M'\) and \(M_*\) are, in the most general case, three independent mass scales. The superfields \(\hat{N}_i, \hat{\phi}\) and \(\hat{S}\) contain the following bosonic components, respectively: the singlet sneutrino inflaton \(\hat{N}\) (dropping the family index here and in the following), which is non-zero during inflation; the so-called waterfall field \(\hat{\phi}\), which is held at zero during inflation but which develops a non-zero vacuum expectation value (vev) after inflation; and the singlet field \(\hat{S}\) which is held practically at zero during and after inflation.

The form of \(W\) in Eq. (1) can be understood as follows:

- The first term on the right-hand side of Eq. (1) serves to fix the vev of the waterfall field after inflation and contributes a large vacuum energy to the potential during inflation.
- The second term on the right-hand side of Eq. (1) allows the sneutrino inflaton to give a positive mass squared for the waterfall field during inflation, which fixes its vev at zero as long as |\(\hat{N}\)| is above a critical value. After inflation, when the waterfall field acquires its non-zero vev, the same term yields the masses of the singlet (s)neutrinos, which are required for explaining the smallness of neutrino masses via the see-saw mechanism.

In order to obtain this minimal form of the superpotential, we have chosen the waterfall superfield to appear in this term as \(\hat{\phi}^4/M'^2\) (instead of e.g. \(\hat{\phi}^2\)) in this example. This allows to impose a \(Z_4\) discrete symmetry which prevents explicit singlet (s)neutrino masses. We would like to note at this point that the dots in Eq. (1) include \(Z_4\)-violating higher-dimensional operators such as, e.g., \(\hat{S}\hat{\phi}^5/M'^3\) (or even \(\hat{S}\hat{\phi}^5/m^3_P\)) which lift the degeneracy of the true vacuum and effectively blow away potential domain wall networks associated with \(Z_4\)-breaking after inflation. \(W\) is also compatible with a \(U(1)_R\)-symmetry under which \(W\) and \(\hat{S}\) each carry unit R-charge, while the charge of \(\hat{N}\) is 1/2. Under suitable conditions the discrete subgroup of this symmetry acts as matter parity.

THE KÄHLER POTENTIAL

Since the field values of the inflaton are well below the reduced Planck scale \(m_P = 1/\sqrt{8\pi G_N}\), we can consider an expansion in powers of \(1/m_P^2\):

$$\mathcal{K} = |\hat{S}|^2 + |\hat{\phi}|^2 + |\hat{N}|^2 + \kappa_S \frac{|\hat{S}|^4}{4m_P^2} + \kappa_N \frac{|\hat{N}|^4}{4m_P^2} + \kappa_\phi \frac{|\hat{\phi}|^4}{4m_P^2}$$

$$+ \kappa_{\hat{S}\hat{\phi}} \frac{|\hat{S}|^2|\hat{\phi}|^2}{m_P^2} + \kappa_{\hat{S}\hat{N}} \frac{|\hat{S}|^2|\hat{N}|^2}{m_P^2} + \kappa_{\hat{N}\hat{\phi}} \frac{|\hat{N}|^2|\hat{\phi}|^2}{m_P^2} + \ldots.$$  \hspace{1cm} (2)
where the dots indicate higher order terms and additional terms for the other fields.

With non-zero F-terms during inflation, the non-canonical Kähler potential can contribute significantly to the scalar potential.

**THE SCALAR POTENTIAL**

Within the model defined by the superpotential $\mathcal{W}$ of Eq. (1) and the Kähler potential $\mathcal{K}$ of Eq. (2), we can now analyze the scalar potential. The F-term contributions to the scalar potential are given by:

$$V_F = e^{\mathcal{K}/m_P^2} \left[ K^{-1}_{ij} D_{z_i} \mathcal{W} D_{z_j} \mathcal{W}^* - 3m_P^2 |\mathcal{W}|^2 \right],$$

(3)

with $z_i$ being the bosonic components of the superfields $\hat{z}_i \in \{ \hat{N}, \hat{\phi}, \hat{S}, \ldots \}$ and where we have replaced the superfields in $\mathcal{W}$ and $\mathcal{K}$ by their bosonic components and defined

$$D_{z_i} \mathcal{W} := \frac{\partial \mathcal{W}}{\partial z_i}, \quad K_{ij} := \frac{\partial^2 \mathcal{K}}{\partial z_i \partial z_j^*}$$

(4)

and $D_{z_j} \mathcal{W}^* := (D_{z_j} \mathcal{W})^*$. Since we assume that $\hat{N}, \hat{\phi}$ and $S$ are effective gauge singlets at the energy scales under consideration, there are no relevant D-term contributions. From Eqs. (1) and (2), with canonically normalized fields, and writing the potential in terms of real fields $\hat{N}_R = \sqrt{2} |\hat{N}|$, $\phi_R = \sqrt{2} |\phi|$ and $S_R = \sqrt{2} |S|$, we obtain

$$V = \kappa^2 \left( \frac{\phi_R^4}{4M^2} - M^2 \right)^2 \left( 1 - \beta \frac{\phi_R^2}{2m_P^2} + \gamma \frac{\tilde{N}_R^2}{2m_P^2} - \kappa_S \frac{S_R^2}{2m_P^2} \right)$$

$$+ \frac{\gamma^2}{2M^4} (\tilde{N}_R^2 \phi_R^2 + \tilde{N}_R^4 \phi_R^4) + \ldots,$$

(5)

where we have defined

$$\beta := \kappa_S \phi - 1 \quad (> 0 \text{ for inflation to end}),$$

$$\gamma := 1 - \kappa_{SN}.$$  

(6)

(7)

**REALIZING SNEUTRINO HYBRID INFLATION**

From Eq. (5) we see that $S_R$ can be set to zero during inflation if we take, e.g., $\kappa_S < -1/3$, such that $S_R$ gets a mass term larger than the Hubble parameter $H \approx \sqrt{V_0}/(\sqrt{3}m_P)$. With $\phi_R = S_R = 0$, the part of the scalar potential relevant for the evolution of the singlet sneutrino inflaton $\hat{N}_R$ during inflation is given by

$$V = \kappa^2 M^4 \left( 1 + \gamma \frac{\tilde{N}_R^2}{2m_P^2} + \delta \frac{\tilde{N}_R^4}{4m_P^4} \right) + \ldots,$$

(8)

1 We will neglect radiative corrections to the potential in the following, which are generically subdominant in our model.
FIGURE 1. Inflationary trajectory and end of inflation by a second order phase transition in sneutrino hybrid inflation.

where we have included the next-to-leading order term proportional to $\delta$.

During inflation, the waterfall field $\phi_R$ has a zero vev and the potential is dominated by the vacuum energy $V_0 = \kappa^2 M^4$. This false vacuum during inflation is stable as long as the mass squared for the waterfall field $\phi_R$ is positive. From Eq. (5) we obtain the requirement

$$m_{\phi_R}^2 = \lambda N \frac{\bar{N}_R^4}{M_*^2} - \beta \frac{\kappa^2 M^4}{m_P^2} > 0.$$  \hspace{1cm} (9)

Inflation thus ends when the squared mass of the waterfall field becomes negative, i.e. $\phi_R$ develops a tachyonic instability and rolls rapidly to its global minimum at $\langle \phi_R \rangle = (2M M')^{1/2}$, ending inflation (as illustrated in Fig. 1). Clearly, this requires $\beta > 0$, as already indicated in Eq. (6). More precisely, inflation ends by a second order phase transition when the field value of the inflaton drops below the critical value $\bar{N}_R$ given by

$$\bar{N}_R^2 = \sqrt{\frac{\kappa M^2 M_*}{\lambda N}} \frac{m_P}{m_P}.$$  \hspace{1cm} (10)

The field value of the inflaton $\bar{N}_R$ at $N = 50$ to 70 $e$-folds before the end of inflation is then given approximately by

$$\bar{N}_R \approx \bar{N}_R e^{\gamma N}.$$  \hspace{1cm} (11)

During inflation, the parameter $\gamma$ in the scalar potential in Eq. (8) controls the mass of the inflaton. Furthermore, compared to the term proportional to $\gamma$, the term proportional to $\delta$ is suppressed by $\bar{N}_R^2/m_P^2$. The slow-roll parameters are given by

$$\varepsilon := \frac{m_P^2}{2} \left( \frac{V'}{V} \right)^2 \approx \frac{(\delta \bar{N}_R^3 + m_P^2 \gamma \bar{N}_R)^2}{2m_P^2} \approx \gamma^2 \frac{\bar{N}_R^2}{2m_P^2},$$  \hspace{1cm} (12)

$$\eta := m_P^2 \left( \frac{V''}{V} \right) \approx \gamma + \frac{3 \delta \bar{N}_R^2}{m_P^2} \approx \gamma,$$  \hspace{1cm} (13)

$$\xi := m_P^4 \left( \frac{V'''}{V^2} \right) \approx 6 \frac{\delta \bar{N}_R^2 (\gamma m_P^2 + \delta \bar{N}_R^2)}{m_P^4}.$$  \hspace{1cm} (14)
FIGURE 2. Graphical illustration of the predictions for the tensor-to-scalar ratio $r$ and the spectral index $n_s$ in sneutrino hybrid inflation [5] (narrow strip on the bottom) and in chaotic sneutrino inflation [3] with a quadratic potential (white dots), compared to the allowed regions by WMAP [7]. In sneutrino hybrid inflation, the tensor-to-scalar ratio $r$ is predicted to be $r \ll \gamma^2$, with $\gamma = 0.025 \pm 0.01$ from the experimental data on the spectral index, as discussed in the text. By their predictions for $r$, future observations will be able to distinguish sneutrino hybrid inflation from chaotic sneutrino inflation.

where prime denotes derivative with respect to $N_R$. Thus, assuming that the slow-roll approximation is justified (i.e. $\epsilon \ll 1, \eta \ll 1$), the spectral index $n_s$, the tensor-to-scalar ratio $r = A_t/A_s$ and the running spectral index $dn_s/d\ln k$ are given by

$$n_s \approx 1 - 6\epsilon + 2\eta \approx 1 + 2\gamma,$$  \hspace{1cm} (15)

$$r \approx 16\epsilon \approx \gamma^2 \frac{8\tilde{N}_R^2}{m_P^2},$$  \hspace{1cm} (16)

$$\frac{dn_s}{d\ln k} \approx 16\epsilon\eta - 24\epsilon^2 - 2\xi \approx -\gamma \frac{12\delta\tilde{N}_R^2}{m_P^2}.$$  \hspace{1cm} (17)

CONSTRAINTS FROM EXPERIMENTAL DATA

The experimental data on the spectral index from WMAP $n_s = 0.95 \pm 0.02$ [7] restricts $\gamma$ to be roughly $\gamma = 0.025 \pm 0.01$. As discussed above, $\gamma$ controls the sneutrino mass during inflation. In this model it stems mainly from supergravity corrections. In addition, we see that the tensor-to-scalar ratio $r = A_t/A_s$ and the running spectral index $dn_s/d\ln k$ are suppressed by higher powers of $\gamma$ or by $\tilde{N}_R^2/m_P^2$ and are thus generically small. Especially the prediction for the tensor-to-scalar ratio $r \ll \gamma^2$ is thus in sharp contrast to the prediction of $r \approx 0.16$ for the case of chaotic sneutrino inflation with a quadratic superpotential, as illustrated in Fig. 2.
In our model, the amplitude of the primordial spectrum is given by

\[ P_{\mathcal{R}}^{1/2} \approx \frac{1}{\sqrt{2\varepsilon}} \left( \frac{H}{2\pi m_p} \right) \approx \frac{\kappa}{2\sqrt{3} \sqrt{\gamma \pi m_p} \tilde{N}_{\text{Re}}} \cdot \]  

(18)

Given the COBE normalization \( P_{\mathcal{R}}^{1/2} \approx 5 \times 10^{-5} \), from Eqs. (10), (18) and (11) we obtain

\[ \frac{M^2}{M_* m_p} \approx 3 \times 10^{-8} \frac{\gamma^2 \sqrt{B}}{\kappa \lambda_N}, \]  

(19)

which relates the scale \( M \) in the superpotential to the cutoff scale \( M_* \). It has to be combined with the constraint \( \tilde{N}_{\text{Re}} \ll m_p \) (see Eqs. (11) and (10)) and with \( \Reheating < M' < M_* \).

**REHEATING AND NON-THERMAL LEPTOGENESIS**

An interesting feature of sneutrino inflation is that the observed baryon asymmetry can arise via non-thermal leptogenesis [2, 6] directly through sneutrino inflaton decay [9]. To illustrate the mechanism, let us assume in the following the situation that the inflaton is the lightest singlet sneutrino \( \tilde{N}_1 \) and that it dominates leptogenesis and reheating after inflation. In sneutrino hybrid inflation [5], this is, e.g., the case if the waterfall field \( \phi \) decays earlier than the singlet sneutrino inflaton via heavier singlet neutrinos \( N_2 \) (or \( N_3 \)) with comparably large couplings to \( \phi \).

From Eq. (1), using \( \langle \phi \rangle = \sqrt{M' M} \), we see that the mass of \( \tilde{N}_1 \) is given by \( M_{R1} = 2(\lambda_N)_{11} M' M / M_* \) in the basis where the mass matrix \( M_R \) of the singlet (s)neutrinos is diagonal. It decays mainly via the extended MSSM Yukawa coupling \( (Y_\nu)^i_1 \tilde{L}_i \tilde{H}_u \tilde{N}_1 \) into slepton and Higgs or into lepton and Higgsino with a decay width given by \( \Gamma_{N_1} = M_{R1} (Y_\nu)^i_1 (Y_\nu)^i_1 / (4\pi) \). The decay of the singlet sneutrino after inflation reheats the universe to a temperature \( T_{\Reh} \approx (90 / (228.75 \pi^2))^{1/4} \sqrt{T_{N_1} m_p} \). If \( M_{R1} > T_{\Reh} \), the lepton asymmetry is produced via cold decays of the singlet sneutrinos and the produced baryon asymmetry can be estimated as \( n_B/n_\gamma \approx -1.84 \varepsilon_1 T_{\Reh} / M_{R1} \), where \( \varepsilon_1 \) is the decay asymmetry for the singlet sneutrino \( \tilde{N}_1 \).

To take a concrete example of the above discussion, neutrino Yukawa couplings \( (Y_\nu)^i_1 \approx 10^{-6} \) and a sneutrino mass of \( M_{R1} = 10^8 \) GeV allow to generate the observed baryon asymmetry of the universe \( n_B/n_\gamma \approx (6.1 \pm 0.2) \times 10^{-10} \) [7] and imply a reheat temperature \( T_{\Reh} \approx 10^6 \) GeV. Such a low reheat temperature is desirable with respect to gravitino constraints (see e.g. [10]) in some supergravity models.

**SUMMARY AND CONCLUSIONS**

We have reviewed the scenario of sneutrino hybrid inflation, where the singlet sneutrino, the superpartner of the right-handed neutrino, plays the role of the inflaton. Sneutrinos are present in any extension of the MSSM, where the smallness of the observed neutrino masses is explained via the see-saw mechanism. In a minimal model of sneutrino hybrid inflation in supergravity, we have found a spectral index \( n_s \approx 1 + 2\gamma \) with \( \gamma = 0.025 \pm \)
0.01 constrained by WMAP, leading to the prediction $|dn_s/d\ln k| \ll |\gamma|$ for the running spectral index and $r \ll \gamma^2$ for the tensor-to-scalar ratio. The prediction for the tensor-to-scalar ratio in sneutrino hybrid inflation is thus much smaller than the prediction $r \approx 0.16$ of chaotic sneutrino inflation and makes sneutrino hybrid inflation easily distinguishable from chaotic sneutrino inflation by future observations. In contrast to chaotic inflation, the field values of the singlet sneutrino inflaton in hybrid inflation are well below the Planck scale. We have discussed how the baryon asymmetry of our universe can be explained via non-thermal leptogenesis and how a low reheat temperature $T_{RH} \approx 10^6$ GeV can be realized with neutrino Yukawa couplings consistent with first family quark and lepton Yukawa couplings in Grand Unified Theories.

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