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Sneutrino warm inflation in the minimal supersymmetric model

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The model of RH neutrino fields coupled to the MSSM is shown to yield a large parameter regime of warm inflation. In the strong dissipative regime, it is shown that inflation, driven by a single sneutrino field, occurs with *all field amplitudes below the Planck scale*. Analysis is also made of leptogenesis, neutrino mass generation and gravitino constraints. A new warm inflation scenario is purposed in which one scalar field drives a period of warm inflation and a second field drives a subsequent phase of reheating. Such a model is able to reduce the final temperature after inflation, thus helping to mitigate gravitino constraints.

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I. INTRODUCTION

In recent times, the idea of inflation being driven by the bosonic supersymmetric partner to a neutrino field has generated interest [1,2]. The idea is not new [3,4], but impetus has been gained after the experimental discovery of neutrino masses and mixing and an explanation through the seesaw mechanism [5]. In supersymmetric realizations of the seesaw mechanism, the right-handed neutrinos have bosonic partners, sneutrinos, which are singlet fields, thus possible inflaton candidates. Model building typically proceeds by simply adding on the additional right-handed neutrino fields to an existing model. Thus the simplest supersymmetric model that emerges is an extended version of the MSSM, with now three families of right-handed neutrinos added on.

Two types of sneutrino inflation models have been examined, chaotic [1] and hybrid [2] sneutrino inflation. The chaotic model is the simplest to construct, since all it requires is a monomial potential which can easily be obtained directly from the sneutrino fields. However this model suffers from the large field problem, in that the sneutrino field that drives inflation will have to have a field amplitude above the Planck scale. In the effective field theory interpretation of global Supersymmetric models, they are regarded as low-energy limits of some more complete supergravity (sugra) theory. However, for example, in ''minimal'' sugra the exponential factor in front of the potential would prevent any scalar field from getting a value larger than m_P. Chaotic inflation would be possible with other more involved choices of the Khaler potential [4] such that sugra corrections are kept under control. Still, in general in these models there are an infinite number of nonrenormalizable operators suppressed by the Planck scale. As such, once the field amplitude exceeds this scale, an infinite number of parameters would require fine-tuning, so leaving no predictability in the theory. It is for this reason that chaotic inflation models are not amenable to particle physics model building. Hybrid inflation scenarios overcome the large field problem, since all field amplitudes are well below the Planck scale. However for sneutrino inflation, these models require introducing two additional superfields aside from the right-handed neutrino fields [2]. As such, this model is more contrived than the chaotic model. Nevertheless, up to now the hybrid model appears to be the simplest model in which to implement sneutrino inflation and be amenable to particle physics model building.

In this paper an even simpler model of sneutrino inflation is presented. In particular we show that monomial potentials, which can be constructed with only the right-handed sneutrino fields, when coupled to the MSSM, realize warm inflationary regimes. We show that in such regimes, due to the effect of strong dissipation, the field amplitudes of all sneutrino fields are well below the Planck scale, thus allowing such models to be consistent with particle physics model building.

The paper is organized as follows. The basic model is presented in Sec. II. The dissipative effects and basic equations of warm inflation for this model are obtained in Sect. III. The results of the sneutrino warm inflation scenario, which incorporates leptogenesis, are given in Sec. IV. An issue that emerges in Sect. IV is that the final temperature after inflation is too large to adequately control gravitino constraints. To improve this situation, in Sec. V a new warm inflation scenario is presented. In Sec. 4 neutrino mass generation from this scenario are examined. Finally in Sec. VII we summarize our results.

II. MODEL

We consider the model of three generations of right-handed neutrinos, N_i, coupled to the MSSM with the

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superpotential

$$W = \frac{M_{Ni}}{2} N_i N_i + (h_N)_{ij} H_u L_i N_j + (h_L)_{ij} H_d L_i E_j^c + (h_u)_{ij} H_u Q_i U_j^c + (h_d)_{ij} H_d Q_i D_j^c. \qquad (1)$$

The above model contains all the usual MSSM matter superfields, H_u, H_d the Higgs doublets giving masses to the up and down quarks, respectively, Q_i, the left-handed quarks, U_i, D_i, the right-handed up and down quarks, respectively, and L_i, E_i the left-handed and right-handed leptons. During inflation, assuming that at least one of the sneutrinos has got a nonzero vacuum expectation value (vev), the relevant terms in the potential are:

$$V = M_{Ni}^2 |N_i|^2 + |(h_N)_{ij}|^2 |L_i H_u|^2 + 2Re[(h_N^*)_{ij} M_i N_j L_i^* H_u^*] + |(h_N)_{ij}|^2 (|H_u|^2 + |L_i|^2)|N_j|^2 + \ldots . \qquad (2)$$

For a large value of N_i we do not have to worry about soft SUSY breaking terms, and then all the spectrum remains massless ($H_u = H_d = 0$), except for the fields that couple directly to the sneutrinos, i.e., H_u and the lepton doublets L_j. With $\langle N_i \rangle = \phi_{Ni}/\sqrt{2}$, the scalars, for example, get masses

$$m_{H_u}^2 = \frac{1}{2} \sum_j h_{Nj}^2 \phi_{Nj}^2 \qquad (3)$$

and similarly for the sleptons, where

$$h_{N_j}^2 = \sum_i |(h_N)_{ij}|^2. \qquad (4)$$

III. DISSIPATIVE INFLATIONARY DYNAMICS

The interaction of the inflaton with other fields leads in general not only to modifications of the inflaton effective potential, but also to dissipative effects [6,7]. These effects result in radiation production during inflation as well as modify the inflaton evolution equation with energy nonconserving terms. If these dissipative effects are adequately large, they can alter the standard picture of inflation, leading to warm inflation [8].

An analysis of various interaction configuration [6,7] has shown that warm inflation occurs generically in many typical inflaton models. For example recently we showed in [9] that the popular SUSY hybrid inflation model has a sizable parameter regime of warm inflation. In this section we show that warm inflation occurs in the sneutrino-MSSM model Eq. (1).

A basic interaction structure that has been shown in [6,7,10] to produce sizable dissipative effects has the form of the bosonic inflaton field coupled to a heavy bosonic field which in turn is coupled to a light fermionic field. Such a structure can easily be identified in the sneutrino-MSSM model Eq. (1). For this consider the simplest case where only one sneutrino dominates the energy density during inflation, say N_1, thus acting the role of the inflaton field. Then from Eq. (1) the following relevant interaction configuration can be extracted

$$\mathcal{L}_I = -|h_N|^2 |N_1|^2 |H_u|^2 + h_t H_u \bar{t}_R t_L + h.c., \qquad (5)$$

thus the inflation N_1 couples to the up Higgs field, and for a large amplitude for N_1, the H_u field then becomes heavy. This Higgs field in turn is coupled to the top fermion fields, which are massless during inflation. Dissipative effects occur because as the inflaton amplitude changes, it implies a change to the H_u mass. This results in a coherent excitation of the H_u field, which then decays into the light top fermions with decay rate

$$\Gamma_t = \frac{3h_t^2}{16\pi} m_{H_u}. \qquad (6)$$

From the dissipative calculations in Refs. [6,7] this sort of interaction leads to the effective inflaton evolution equation

$$\ddot{\phi}_N + (3H + \Upsilon_N)\dot{\phi}_N + V' = 0, \qquad (7)$$

where $\dot{\phi}_N$ ($\ddot{\phi}_N$) is the first (second) time derivative of the field, and a "prime" denotes the derivative of the potential with respect to the inflaton field. The dissipative coefficient, based on the results in [6,7], can be determined to be

$$\Upsilon_N \simeq \frac{\sqrt{\pi}}{20} Y_N^{3/2} Y_t \phi_N, \qquad (8)$$

with $Y_N \equiv h_N^2/(4\pi)$ and $Y_t = h_t^2/4\pi$. Also in Eq. (7) the potential and the Hubble parameter are

$$V \simeq \frac{1}{2} M_N^2 \phi_N^2 \qquad (9)$$

and

$$H^2 \simeq \frac{M_N^2 \phi_N^2}{6m_P^2}. \qquad (10)$$

The dissipative term in Eq. (7) leads to radiation production which in the expanding spacetime obeys the equation

$$\dot{\rho}_R + 4H\rho_R = \Upsilon_N \dot{\phi}_N^2. \qquad (11)$$

Although the basic idea of interactions leading to dissipative effects during inflation is generally valid, the above set of equations has strictly been derived in [6,7] only in the adiabatic-Markovian limit, i.e., when the fields involved are moving slowly, which requires

$$\frac{\dot{\phi}_N}{\phi_N} < H < \Gamma_t, \qquad (12)$$

with Γ_t being the decay rate Eq. (6). The second inequality, $H < \Gamma_t$ is also the condition for the radiation (decay products) to thermalize.

Thus in general any inflation model could have two very distinct types of inflationary dynamics, which have been termed cold and warm [6–8]. The cold inflationary regime is synonymous with the standard inflation picture [11–13], in which dissipative effects are completely ignored during the inflation period. On the other hand, in the warm inflationary regime dissipative effects play a significant role in the dynamics of the system. A rough quantitative measure that divides these two regimes is $\rho_R^{1/4} \approx H$, where $\rho_R^{1/4} > H$ is the warm inflation regime and $\rho_R^{1/4} \lesssim H$ is the cold inflation regime. This criteria is independent of thermalization, but if such were to occur, one sees this criteria basically amounts to the warm inflation regime corresponding to when $T > H$. This is easy to understand since the typical inflaton mass during inflation is $m_\phi \approx H$ and so when $T > H$, thermal fluctuations of the inflaton field will become important. This criteria for entering the warm inflation regime turns out to require the dissipation of a very tiny fraction of the inflaton vacuum energy during inflation. For example, for inflation with vacuum (i.e. potential) energy at the GUT scale $\sim 10^{15-16}$ GeV, in order to produce radiation at the scale of the Hubble parameter, which is $\approx 10^{10-11}$ GeV, it just requires dissipating one part in 10^{20} of this vacuum energy density into radiation. Thus energetically not a very significant amount of radiation production is required to move into the warm inflation regime. In fact the levels are so small, and their eventual effects on density perturbations and inflaton evolution are so significant, that care must be taken to account for these effects in the analysis of any inflation models.

The conditions for slow-roll inflation ($\dot\phi_N^2 \ll V$, $\ddot\phi_N \ll H\dot\phi_N$) are modified in the presence of the extra friction term Υ_N, and the slow-roll parameters are given now by:

$$\epsilon_\Upsilon = \frac{\epsilon_H}{(1+r)^2}, \tag{13}$$

$$\eta_\Upsilon = \frac{\eta_H}{(1+r)^2}, \tag{14}$$

where

$$r \equiv \frac{\Upsilon_N}{3H} = \frac{\sqrt{6\pi}}{60} Y_N^{3/2} Y_t \frac{m_P}{M_N} = C_\Upsilon \frac{m_P}{M_N}. \tag{15}$$

and $\epsilon_H \equiv m_P^2 V'^2/(2V^2)$, $\eta_H \equiv m_P^2 V''/V$ are the standard cold inflation slow-roll parameters, in which there are no dissipation effects. Therefore, when $\Upsilon_N \gg 3H$ ($r \gg 1$), we have:

$$\epsilon_\Upsilon \simeq \frac{2}{C_\Upsilon^2}\left(\frac{M_N}{\phi_N}\right)^2, \tag{16}$$

$$\eta_\Upsilon \simeq \frac{2}{C_\Upsilon^2}\left(\frac{M_N}{\phi_N}\right)^2. \tag{17}$$

In the slow-roll regime, when $\eta_\Upsilon < 1$ and $\epsilon_\Upsilon < 1$, Eqs. (7) and (11) are well approximated by:

$$\dot\phi_N \simeq -\frac{V'}{3H}\frac{1}{1+r}, \tag{18}$$

$$\rho_R \simeq \frac{\Upsilon_N}{4H}\dot\phi_N^2 \simeq \frac{1}{2}\frac{r}{(1+r)^2}\epsilon_H V, \tag{19}$$

and the number of e-folds is given by:

$$N_e \simeq -\int_{\phi_{\rm end}}^{\phi_{Ne}} \frac{3H^2}{V'}(1+r)d\phi_N \simeq \frac{(1+r)}{4m_P^2}(\phi_{Ne}^2 - \phi_{\rm end}^2), \tag{20}$$

where $\phi_{Ne}(\phi_{end})$ is the value of the field at 60 e-folds (end of inflation). Inflation ends when $\epsilon_\Upsilon \simeq \eta_\Upsilon \simeq 1$ or when $\rho_R \simeq \rho_N$, whatever happens first. In the former case we have $\phi_{\rm end}^2 \simeq 2m_P^2/(1+r)^2$, whereas for $\rho_R \simeq \rho_N$ then $\phi_{\rm end}^2 \simeq m_P^2 r/(1+r)^2$. In either case, taking $r \gg 1$, we get

$$\phi_{Ne} \simeq \sqrt{\frac{4N_e}{1+r}} m_P. \tag{21}$$

If we also want to keep the field below the Planck scale, we need $r > 4N_e \simeq 240$. From Eq. (15), taking $|h_N| \simeq 1$, this gives the upper bound on the sneutrino mass $M_N \lesssim 8 \times 10^{12}$ GeV.

The effect of the dissipative term, in addition to producing a friction term for the inflaton field, leads to radiation production which can alter density perturbations. Approximately, one can say that when the radiation production leads to $T > H$, the fluctuations of the inflaton field are induced by the thermal fluctuations, instead of being vacuum fluctuations, with a spectrum proportional to the temperature of the thermal bath. We notice that having $T > H$ does not necessarily require $\Upsilon_N > 3H$. Dissipation may not be strong enough to alter the dynamics of the background inflaton field, but it can be enough even in the weak regime to affect its fluctuations, and therefore the spectrum. Depending on the different regimes, the spectrum of the inflaton fluctuations $P_{\delta\phi}^{1/2}$ is given for cold inflation [14], weak dissipative warm inflation [15,16], and strong dissipative warm inflation [17] respectively by

$$T < H: P_{\delta\phi}^{1/2}|_{T=0} \simeq \frac{H}{2\pi}, \tag{22}$$

$$\Upsilon_N < H < T: P_{\delta\phi}^{1/2}|_T \simeq \sqrt{TH} \sim \sqrt{\frac{T}{H}} P_{\delta\phi}^{1/2}|_{T=0}, \tag{23}$$

$$\Upsilon_N > H: P_{\delta\phi}^{1/2}\Big|_\Upsilon \simeq \left(\frac{\pi\Upsilon_N}{4H}\right)^{1/4}\sqrt{TH} \sim \left(\frac{\pi\Upsilon_N}{4H}\right)^{1/4} \times \sqrt{\frac{T}{H}} P_{\delta\phi}^{1/2}\Big|_{T=0}, \tag{24}$$

with the amplitude of the primordial spectrum of the curvature perturbation given by:

$$P_{\mathcal{R}}^{1/2} = \left|\frac{H}{\dot{\phi}_S}\right| P_{\delta\phi}^{1/2} \simeq \left|\frac{3H^2}{V'}\right|(1+r)P_{\delta\phi}^{1/2}. \tag{25}$$

Given the different "thermal" origin of the spectrum, the spectral index also changes with respect to the cold inflationary scenario [18–21], even in the weak dissipative warm inflation regime when the evolution of the inflaton field is practically unchanged. General expressions for the spectral index are given in [9] and those relevant to the model in this paper will be given in the sections that follow.

IV. SNEUTRINO WARM INFLATION

As we have seen, depending on the value of the dissipative coefficient Υ_N, and therefore that of the ratio $r = \Upsilon_N/(3H)$, we can have standard cold inflation or warm inflation (with weak or strong dissipation). In sneutrino inflation with the minimal matter content of the MSSM plus 3 generations of RH (s)neutrinos as shown in Eq. (1), there is a well define dissipation channel during inflation due to the coupling of the RH sneutrinos to the Higgs H_u, and the coupling of H_u to the top sector. From the experimental value of the top quark mass, $m_t = 174(178)$ GeV [22], the value of the top Yukawa coupling h_t at the electroweak scale has to be close to 1, with $m_t \simeq h_t(m_t)v\sin\beta$, $v = \sqrt{v_u^2 + v_d^2} = 174$ GeV being the Higgs vacuum expectation value (vev), and $\tan\beta = \langle H_u\rangle/\langle H_d\rangle$ the ratio of the Higgs vevs. The top Yukawa coupling increases due to the running with the scale, and depending on the value of $\tan\beta$ it can reach the perturbative bound $h_t(M_X) \simeq \sqrt{4\pi}$ at the unification scale M_X. Thus, although slightly model dependent, its value at the inflationary scale will be in the range $[1, \sqrt{4\pi}]$. Then, without lost of generality we can take its value close to the perturbative bound $Y_t = h_t^2/(4\pi) \simeq 1$, so that the dissipative coefficient Eq. (8) becomes $\Upsilon_N \approx \frac{\sqrt{\pi}}{20} Y_N^{3/2}\phi_N$. The free parameters in the model are then the sneutrino-inflaton mass M_N and its Yukawa coupling $|h_N^2|$ defined in Eq. (4). Imposing the COBE normalization on the amplitude of the primordial spectrum of perturbations [23,24] generated during inflation, we can fix one of these parameters, say the mass M_N, as dependent on the value of the coupling h_N. This is plotted in Fig. (1), where we have included in addition to the value of M_N (GeV) the values of the dissipative ratio $r = \Upsilon_N/(3H)$, the field value in m_P units at 60 e-folds, the temperature of the thermal bath at the end of inflation T_{end}, and the ratio T/H during inflation. We can clearly distinguish the different regimes in the plot depending on the sneutrino Yukawa value. For very small values $|h_N| < 10^{-3}$ we recover the standard cold inflation predictions, with $\phi_N > m_P$ and $M_N \simeq 2\times 10^{13}$ GeV. In this regime, the Yukawa coupling plays no role during inflation, and the normalization of the spectrum is set by the RH sneutrino mass, with $M_N \simeq 2\times 10^{13}$ GeV [3]. The value of the Yukawa coupling fixes the decay rate of the sneutrino and therefore the final reheating T. In the simplest scenario where the inflaton is the lightest RH sneutrino, we would need $|h_N| < O(10^{-7} - 10^{-6})$ if we want to keep $T_{RH} \leq 10^7 - 10^8$ GeV in order not to have problems with thermal production of gravitinos [25]. We remark that $T_{\rm end}$ in Fig. 1 is the temperature associated to the radiation energy density at the end of inflation due to dissipative effects, but this is not necessarily the reheating T typically defined as the T at which the inflaton completely decays and the Universe becomes radiation dominated. In the cold inflation scenario the radiation energy density at the end of inflation is always subdominant, and then reheating would proceed as usual by the subsequent decay of the sneutrino.

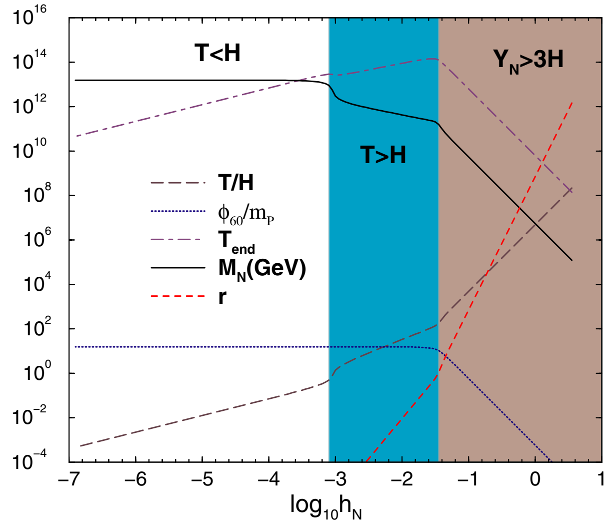

FIG. 1 (color online). Values of the parameters M_N, r and ϕ_{60} with respect to the sneutrino Yukawa coupling $|h_N|$, for cold inflation ($T < H$), weak dissipative regime ($T > H$, $\Upsilon_N < 3H$), and strong dissipative regime (T, $\Upsilon_N/3 > H$).

On the other hand, for a coupling $10^{-3} < |h_N| < 10^{-1}$ inflation takes place in the weak dissipative regime, which would require a sneutrino mass of the order of $10^{12} - 10^{11}$ GeV in order to fit the COBE amplitude of the spectrum. We notice that for these coupling and mass values, $|h_N| > 10^{-3}$ and $M_N < 10^{13}$ GeV, the adiabatic-Markovian approximation Eq. (12) holds. Still, the field values are larger than m_P. Again, the energy density in radiation is not dominant at the end of inflation, and $T_{\rm end}$ is not necessarily the final T_{RH}. However, given that now we have a larger value of the coupling $|h_N|$, the standard estimation of $T_{RH} = (90/(\pi^2 g_*))^{1/4}\sqrt{\Gamma_N m_P}$ in this regime would give a value $O(10^{12}$ GeV), beyond the gravitino constraint[1].

[1]This bound does not apply if gravitinos are the lightest stable SUSY particles [26], like in gauge mediated susy breaking models [27,28]. Also the constraint can be relaxed for very massive gravitino, like, for example, in anomaly mediated susy breaking models [29].

More interesting is the strong dissipative regime with $r > O(100)$ for $|h_N| \gtrsim 0.1$. In this regime field values are always kept below the cut-off scale m_P, which render the theory more attractive from the point of view of particle physics. The model can be considered as an effective model valid below the cut-off scale m_P, without the need of worrying about sugra corrections. Those are kept negligible for field values below the Planck scale. In addition, even with the choice of minimal Kahler potential the sugra model can fulfil the conditions for slow-roll inflation and it does not suffer from the so-called "eta problem" [30]. Even if the exponential factor in the sugra potential gives rise to a mass term contribution of the order of the Hubble parameter for the sneutrino, the friction term in the evolution of the inflaton field is controlled now by the dissipative coefficient with $\Upsilon_N > 3H$, and the slow-roll parameter parameter is given by $\eta_\Upsilon \simeq 3V''/\Upsilon_N^2$. The role of the Hubble expansion rate in the slow-roll conditions is now played by Υ_N, and slow-roll inflation is still possible for masses larger than H but smaller than Υ_N.

More specifically, with minimal Kahler potential for the sneutrino field, $K = \phi_N^2/2$, the sugra potential during inflation (setting the Higgs and slepton fields to zero) is given by:

$$V = e^K(|K_N W + W_N|^2 - 3|W|^2) = e^{\phi_N^2/(2m_P^2)} \frac{M_N^2}{2} \phi_N^2 \left(1 + \frac{\phi_N^2}{8m_P^4} + \frac{\phi_N^4}{16m_P^4}\right), \tag{26}$$

and the exponential factor in front of the sugra potential generates a mass term of the order of the Hubble parameter and other contributions coming from the higher order terms in the potential, i.e.,

$$V'' = 3H^2 + M_N^2 e^{\phi_N^2/(2m_P^2)}\left(1 + \frac{11}{4}\frac{\phi_N^2}{m_P^2} + \frac{31}{16}\frac{\phi_N^4}{m_P^4} + \frac{\phi_N^6}{2m_P^6} + \frac{\phi_N^8}{32m_P^8}\right). \tag{27}$$

In the strong dissipative regime we have that the amplitude of the inflaton field is below the Planck scale, $\phi_N < m_P$, so that we can neglect the high order contributions in the potential and derivatives,

$$V'' \simeq 3H^2 + M_N^2 + \cdots \simeq M_N^2\left(1 + \frac{\phi_N^2}{2m_P^2} + \cdots\right) \simeq M_N^2, \tag{28}$$

and the slow-roll parameters reduce to those given in Eqs. (16) and (17).

In this regime, the sneutrino mass value varies between 10^{10} GeV and 2×10^5 GeV, decreasing with the value of the coupling, as seen in Fig. 1. In particular, using Eqs. (24) and (25), the amplitude of the spectrum of primordial curvature perturbation is given by,

$$P_{\mathcal{R}}^{1/2} \simeq \frac{1}{2}\left(\frac{15}{64g_*}\right)^{1/8}(4N_e)^{3/4} C_\Upsilon^{3/8}\left(\frac{M_N}{m_P}\right)^{3/8}, \tag{29}$$

and using $P_{\mathcal{R}}^{1/2} = 5 \times 10^{-5}$, (and $g_* = 228.75$, the number of effective degrees of freedom for the MSSM) we have for $|h_N| \gtrsim 0.1$

$$|h_N| \simeq 8.2 \times 10^{-2}\left(\frac{10^{10}\ \mathrm{GeV}}{M_N}\right)^{1/3}. \tag{30}$$

This equation summarizes the constraint on the coupling and the sneutrino mass in order to have the strong dissipative regime[2]. The larger the coupling $|h_N|$ is, the lighter the RH sneutrino.

In this regime inflation ends when the energy density in radiation becomes comparable to that of the sneutrino field and the Universe becomes radiation dominated. Therefore, in this case $T_{\mathrm{end}} \simeq T_{RH}$, with values that are still larger than the gravitino bound. Another question is about leptogenesis in this scenario. One of the nice and more appealing features of sneutrino inflation is the possibility of relating in principle different pieces of physics like inflation, and neutrino masses and leptogenesis, through the physics of the RH neutrinos and their couplings. The lepton asymmetry $Y_L = n_L/s$ is generated by the out-of-equilibrium decay of the RH sneutrinos, and then reprocessed to the $B - L$ asymmetry by sphaleron processes at a temperature around $T \sim 100$ GeV, generating the observed baryon asymmetry $Y_B = n_B/s \simeq (8.7 \pm 0.4) \times 10^{-11}$ [24]. Successful thermal leptogenesis, with the initial RH sneutrino abundance produced out of the thermal bath, requires [31] $T_{RH} \gtrsim 2 \times 10^9$ GeV, which is fulfilled for $|h_N| \lesssim 1.8$. It also requires a similar bound for the sneutrino mass, $M_N > 2 \times 10^9$ GeV, although the sneutrino dominating during inflation need not necessarily be the one originating the lepton asymmetry. We could have, for example, the lighter one with the larger Yukawa coupling as the inflaton, and the next-to-lightest being responsible for Y_L. Nevertheless, there are models where these bound can be evaded [32].

Before closing this section, we comment on the predictions for the spectral index n_S, and the running of the spectral index $dn_S/d\ln k$, of the primordial spectrum. Those do not vary significantly from one regime to another. In the case of standard cold sneutrino inflation, we have $n_S - 1 \simeq -4/(2N_e + 1) \simeq -0.03$ and $dn_S/d\ln k \simeq -32/(4N_e + 2)^2 \simeq -5 \times 10^{-4}$, whereas in the strong dissipative regime we have $n_S - 1 \simeq -3/(2N_e) \simeq -0.025$ and $dn_S/d\ln k \simeq -26/(4N_e)^2 \simeq -4.5 \times 10^{-4}$. The distinctive prediction comes from the tensor-to-scalar ratio $r_T = P_T/P_{\mathcal{R}}$, with the primordial spectrum of the tensor modes being $P_T \simeq 2(H/2\pi m_P)^2$. Whereas in cold sneu-

[2]The value of $|h_N|$ in Eq. (30) depends on the value of the top Yukawa coupling as $|h_N| \propto Y_t^{-1/3}$. For example, for $Y_t \simeq 1/\sqrt{4\pi}$ we get $|h_N| \simeq 0.2$ for $M_N = 10^{10}$ GeV.

trino inflation, given that the field is larger than m_P, we have [1,3] $r_T \simeq 16\epsilon \simeq 0.16$, in the strong dissipative regime that ratio is highly suppressed, with

$$r_T \simeq 0.22\left(\frac{0.01}{|h_N|}\right)^{12}. \tag{31}$$

Future CMB experiment like Planck [33], and also gravitational wave detectors currently under study [34], are expected to reach a sensitivity for r_T below 0.01. Therefore, the lack of a signal for the primordial spectrum of gravitational waves in future experiments will rule out sneutrino inflation in its more standard version, but not warm sneutrino inflation.

V. LOWERING THE POST-INFLATION TEMPERATURE

Having a not too low reheating temperature at the end of inflation may reintroduced the problem of unwanted relic particles in the model, like gravitinos [25] and moduli fields in sugra theories [35], which can be thermally produced at the end of inflation. Given that they are in general, relatively long-lived light particles, with only gravitational interactions with other matter fields, they can decay by the time of nucleosynthesis and invalidate its predictions, or come to dominate the energy density in the Universe, overclosing it. Such problems may be avoided with a low enough reheating temperature (''low enough'' depending on the mass of the particle and its decay rate), or by diluting its thermal abundance afterwards by entropy production. Depending on the source for the entropy released, different solutions to the gravitino and moduli problem can be found in the literature, like, for example, a period of thermal inflation [36–39], the decay of a modulus field with the appropiated mass value [40], or the decay of a network of domain walls [41] formed at a phase transition after inflation, among others. In each case, it is neccessary to check that the entropy release does not dilute at the same time the baryon asymmetry below observational bounds, or if this is case, that it can be generated afterwards.

In the context of sneutrino inflation, maybe the simpler alternative is considering the decay of a *lighter* sneutrino field other than the inflaton, with a larger decay rate, such that it comes to dominate the energy density before decay [42]. This scenario can be easily implemented at the end of warm inflation with strong dissipation, without the need of having large amplitude values for the fields. Moreover, it is compatible with generating the right level of baryon asymmetry by nonthermal leptogenesis.

Let us denote by ϕ_{N2}, M_{N2} the field and mass parameter of the RH sneutrino dominating the energy density during inflation, and ϕ_{N1}, M_{N1} those of a lighter RH sneutrino. The dominant contribution during inflation is given by ϕ_{N2}, $V \simeq M_{N2}^2\phi_{N2}^2/2$, but still the lighter sneutrino can follow a slow-roll trajectory during inflation for field values below m_P, with the slow-roll parameter for ϕ_{N1} being

$$\eta_1 \simeq 2\frac{M_{N1}^2}{M_{N2}^2}\frac{m_P^2}{\phi_{N2}^2}, \tag{32}$$

$$\epsilon_1 \simeq \eta_1\left(\frac{\phi_{N1}}{\phi_{N2}}\right)^2. \tag{33}$$

In order to have $\eta_1 < 1$, $\epsilon_1 < 1$, we only need to assume that the field values during inflation are comparable, $\phi_{N1} \simeq \phi_{N2} < m_P$, and require $M_{N1}/M_{N2} < (\phi_{N2}/m_P)_{\rm end} \simeq 1/\sqrt{2r}$, which in terms of the coupling reads:

$$\frac{M_{N1}}{M_{N2}} \lesssim 1.6\times 10^{-5}|h_{N2}|^{-3}. \tag{34}$$

Having a second slow-rolling field does not change the primordial spectrum during or after inflation, so the estimation given in Eq. (29) applies, and the spectrum is dominated by thermal effects. Moreover, it does not matter what is the amplitude of the curvature perturbation generated by ϕ_{N1} during inflation, by the end of inflation it has leveled to that of ϕ_{N2}. The constraint on the sneutrino mass dominating during inflation, $M_{N2} \lesssim 10^{10}$ GeV, obtained in the previous section still applies.

During inflation the lightest sneutrino ϕ_{N1} energy density is subdominant. When inflation ends for ϕ_{N2} it does so for ϕ_{N1}. Then, this field, weakly coupled, starts oscillating and its energy density on average behaves like matter. Therefore, if its decay rate is small enough, it will end up dominating over the radiation energy density dissipated by ϕ_{N2} during warm inflation. Later the field decays, and it is at this point that we define the final reheating T. Thus, the inflationary period is controlled by ϕ_{N2}, but the reheating phase is controlled by ϕ_{N1}, with

$$T_{RH} \simeq \left(\frac{90}{\pi^2 g_*}\right)^{1/4}\sqrt{\Gamma_1 m_P}, \tag{35}$$

where $\Gamma_1 \simeq \frac{|h_{N1}|^2}{16\pi}M_{N1}$, and then

$$|h_{N1}| \simeq 0.86\times 10^{-6}\times\left(\frac{T_{RH}}{10^6\ \mathrm{GeV}}\right)\left(\frac{10^8\ \mathrm{GeV}}{M_{N1}}\right). \tag{36}$$

Leptogenesis can now proceed through the out-of-equilibrium decay of the lightest sneutrino ϕ_{N1} during the reheating period. Any previous lepton asymmetry would be diluted by the entropy produced by the ϕ_{N1} decay. The lepton asymmetry at the end of reheating is then given by [3,4],

$$\frac{n_L}{s}|_{RH} \simeq |\epsilon_1|\frac{3T_{RH}}{4M_{N1}}, \tag{37}$$

with ϵ_1 being the *CP* asymmetry generated by the decay, given by the interference of the tree level with the one-loop amplitude,

$$|\epsilon_1| \simeq \frac{3}{8\pi(h_N^\dagger h_N)_{11}}\sum_{i\neq 1} Im[(h_N^\dagger h_N)_{1i}]^2\frac{M_{N1}}{M_i}. \tag{38}$$

The asymmetry parameter is bounded by[3] [31]

$$|\epsilon_1| \le \frac{3}{8\pi}\sqrt{\Delta m_A^2}\frac{M_{N1}}{v_u^2} \simeq 2\times 10^{-8}\left(\frac{M_{N1}}{10^8\ \mathrm{GeV}}\right), \quad (39)$$

where $\Delta m_A^2 \simeq (1.3-4.2)\times 10^{-3}$ eV2 is the the atmospheric neutrino mass squared difference [43,44]. From Eqs. (37) and (39) we have then the bound [1,45],

$$\frac{n_L}{s}\bigg|_{RH} \le |\epsilon_1|\frac{3T_{RH}}{4M_{N1}} \simeq 1.5\times 10^{-10}\frac{T_{RH}}{10^6\ \mathrm{GeV}}, \quad (40)$$

and the baryon asymmetry:

$$\frac{n_B}{s}\bigg|_{RH} \le \frac{8}{23}\frac{n_L}{s}|_{RH} \simeq 5\times 10^{-11}\frac{T_{RH}}{10^6\ \mathrm{GeV}}, \quad (41)$$

and hence $T_{RH} \gtrsim 2\times 10^6$ GeV in order to match the observed baryon asymmetry. On the other hand, in order to ensure the out-of-equilibrium decay of ϕ_{N1} and avoid thermal washout of the asymmetry, we require $M_{N1} \gtrsim T_{RH}$. Using Eqs. (30) and (34), the limiting value [1] $T_{RH} \simeq M_{N1} \simeq 10^6$ GeV is reached for $|h_{N2}| \lesssim 0.24$.

Therefore, we have a narrow window of values $0.1 \lesssim |h_{N2}| \lesssim 0.24$, for which inflation happens in the strong dissipative regime with $M_{N2} \simeq 10^{10} - 10^9$ GeV, but reheating with $T_{RH} \simeq 10^6$ GeV and *nonthermal* leptogenesis is given by the decay of a lighter sneutrino with parameters $5\times 10^8\mathrm{GeV} \gtrsim M_{N1} \gtrsim 10^6$ GeV and $2\times 10^{-7} \lesssim |h_{N1}| \lesssim 8.6\times 10^{-5}$. For having warm inflation with a larger Yukawa coupling $h_{N2} \gtrsim 0.24$, the second lighter and long-lived sneutrino with $M_{N1} < 10^6$ GeV can lower the final reheating T below the gravitino bound, but it does not seem consistent with nonthermal leptogenesis as we cannot satisfy at the same time $M_{N1} \gtrsim T_{RH}$ and $T_{RH} \simeq 10^6$ GeV. It would remain to check whether thermal leptogenesis could be viable during the reheating period in this case, for which one would need to set and study the Boltzmann equations describing the evolution of the different number densities, which is beyond the scope of this paper.

VI. WARM INFLATION AND LIGHT NEUTRINO MASSES

In this section we briefly want to comment on the issue of light neutrino masses with a not too heavy sneutrino $M_N \le 10^{10}$ GeV but large Yukawa couplings $|h_N| \ge 0.1$. Over the recent years, different neutrino experiments have established the existence of neutrino oscillations driven by nonzero neutrino masses and neutrino mixing [46]. Atmospheric neutrino oscillation parameters read:

$$|\Delta m_A^2| = (1.3-4.2)\times 10^{-3}\ \mathrm{eV}^2,$$
$$\sqrt{|\Delta m_A^2|} \simeq 0.05\ \mathrm{eV}, \qquad \sin 2\theta_A \ge 0.85, \quad (42)$$

while solar neutrino oscillation parameters lie in the low-LMA (large mixing angle) solution with:

$$|\Delta m_{\mathrm{sun}}^2| = (8.5-7.4)\times 10^{-5}\ \mathrm{eV}^2,$$
$$\tan\theta_{\mathrm{sun}} = 0.4^{+0.09}_{-0.07}. \quad (43)$$

On the other hand, a combined analysis of the solar neutrino, CHOOZ and KamLAND data gives $\sin^2\theta_{13} < 0.055$. An upper limit on the absolute value of the masses is obtained from WMAP data as $\sum_j m_j < (0.7-2.0)$ eV.

Given the superpotential Eq. (1), light neutrino masses are given by diagonalizing the seesaw mass matrix [5]:

$$m_{LL} = v_u^2 h_N M_{RR}^{-1} h_N^T = U_{\nu L}\mathrm{diag}(m_1, m_2, m_3)U_{\nu L}^T, \quad (44)$$

where[4] $v_u = \langle H_u\rangle \sim 174$ GeV, m_i the light LH neutrino masses, and $U_{\nu L}$ the rotation matrix. In the Yukawa matrix h_N, each column define a vector $\mathbf{h_{Ni}}$ with modulus $|h_{Ni}|$ as given in Eq. (4). In the eigenmass basis for the RH neutrinos we have:

$$Trm_{LL} = \sum_{i=1,2,3} m_i \simeq v_u^2\sum_{i=1,2,3}\frac{\mathbf{h_{Ni}}\cdot\mathbf{h_{Ni}}}{M_{Ni}} < O(1)\ \mathrm{eV}. \quad (45)$$

For the parameters of the strong dissipative regime we clearly exceed the WMAP bound, with $\mathbf{h_{N2}}\cdot\mathbf{h_{N2}}/M_{N2} > O(30)$ eV.

However, this applies when assuming a diagonal mass matrix M_{RR} in Eq. (44). We can work instead with [47] (see also [32])

$$M_{RR} = \begin{pmatrix} 0 & M_{N2} & 0 \\ M_{N2} & 0 & 0 \\ 0 & 0 & M_{N1}\end{pmatrix}, \quad (46)$$

where

$$Trm_{LL} = \sum_{i=1,2,3} m_i \simeq v_u^2\left(\frac{2\mathbf{h_{N3}}\cdot\mathbf{h_{N2}}}{M_{N2}} + \frac{\mathbf{h_{N1}}\cdot\mathbf{h_{N1}}}{M_{N1}}\right) < O(1)\ \mathrm{eV}, \quad (47)$$

such that the large contribution coming from the large Yukawa coupling h_{N2} can be canceled out by choosing an appropriate smaller coupling $h_{N3} \ll h_{N2}$. This kind of scheme gives rise to light neutrino masses with an inverted hierarchy, $m_1^2 \approx m_2^2 \gg m_3^3$. For example, taking $\mathbf{h_{N1}}\cdot\mathbf{h_{N1}}/M_{N1} \ll \mathbf{h_{N3}}\cdot\mathbf{h_{N2}}/M_{N2}$. We have then 2 almost degenerate light neutrino masses, with $|m_1| \simeq |m_2| \simeq \sqrt{|\Delta m_A^2|}$, and a massless one $m_3 \simeq 0$, (small corrections

[3]The bound is given by $|\epsilon_1| \le 3(m_3 - m_1)M_{N1}/(8\pi v^2)$, with $m_{3,1}$ being the light neutrino masses. We have taken for simplicity $m_1 < m_3 \simeq \sqrt{\Delta m_A^2} \simeq O(0.05)$ eV, although we could have for example $m_1 < m_3 \approx O(1)$ eV.

[4]Strictly speaking we have $\langle H_u\rangle = v\sin\beta$, with $v = 174$ GeV and $\tan\beta = \langle H_u\rangle/\langle H_d\rangle$. We are setting $\sin\beta \approx 1$ for order of magnitude estimations.

from the Yukawas h_{N1} gives a nonzero m_3 value). Atmospheric neutrino oscillations are given by oscillations among "13" and "23", while solar data is explained by the oscillation between "12", with $m_1^2 - m_2^2 \simeq \Delta m_{\rm sun}^2$ [47].

The mass parameter M_{N1} can be larger or smaller than M_{N2} as far as light neutrino masses are concerned. Nevertheless, if we choose the hierarchy $M_{N2} > M_{N1}$, the asymmetry parameter corresponding to the decay of the lightest sneutrino is given by:

$$|\epsilon_1| \simeq \frac{3}{16\pi(|h_{N1}|^2)} Im[(\mathbf{h}^*_{\mathbf{N3}} \cdot \mathbf{h}_{\mathbf{N1}})(\mathbf{h}^*_{\mathbf{N2}} \cdot \mathbf{h}_{\mathbf{N1}})] \frac{M_{N1}}{M_{N2}} \simeq \frac{3}{8\pi} \sqrt{|\Delta m_A^2|} \frac{M_{N1}}{v_u^2}, \tag{48}$$

where in the second line we have just assumed that there is no hierarchy among the different components of $\mathbf{h}_{\mathbf{N1}}$ and that phases are such that $Im[]$ is maximal and we saturate the upper bound on the asymmetry parameter.

VII. CONCLUSION

The most important new feature for sneutrino inflation found in this paper is a model in which inflation is driven by just a single sneutrino field that creates a monomial inflationary potential, similar to chaotic sneutrino inflation [1,3,4], but with the key difference that in the model of this paper the inflaton amplitude is below the Planck scale. For particle physics model building, this is an important feature, since this model is then not susceptible to large effects from higher dimensional operators. In particular, it can be embedded in a sugra potential even with minimal Kahler potential for the fields, with the exponential sugra correction in front of the potential remaining small and under control. The chaotic inflation scenario of the cold inflation picture has always been attractive for its simplicity, since it requires just a monomial potential to realize inflation. However the downside of this model for model building has been that the inflaton field amplitude necessary for inflation must be larger than the Planck scale, thus making the model highly susceptible to higher dimensional operator corrections. Now, taking into account dissipation, this simple model with a monomial potential can be regarded as an effective model truly valid below the cut-off scale m_P.

Before the results of this paper, the simplest sneutrino model that could maintain field amplitudes below the Planck scale was a version of hybrid inflation [2]. But this model is more contrived since it requires additional fields aside from the RH neutrino fields. Thus the model of this paper is the simplest realistic realization of sneutrino inflation, in the sense that it is the minimal SUSY extension of the standard model to incorporate supersymmetry and RH neutrinos, it requires no additional fields beyond these to realize inflation, and all higher dimensional operators which inevitably also exist are all suppressed.

The key feature of our model that allowed the inflaton amplitude below the Planck scale with a monomial inflation potential was the presence of large dissipation in the inflaton evolution equation. Moreover as shown in Sec. III, the origin of these dissipative effects arise automatically at a first principles level for this model of RH neutrinos coupled to the MSSM. Thus we believe the model in this paper has several attractive features for building a complete model that is able to describe both particle physics and cosmology.

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