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Affleck–Dine leptogenesis via right-handed sneutrino fields in a supersymmetric hybrid inflation model

Zurab Berezhiani a, b, Anupam Mazumdar c, Abdel Pérez-Lorenzana c, d

a Dipartamento di Fisica, Universitá di L’Aquila, I-67010, Coppito, AQ, and INFN, Laboratori Nazionali del Gran Sasso, I-67010, Assergi, AQ, Italy

b The Andronikashvili Institute of Physics, GE-380077 Tbilisi, Georgia

c The Abdus Salam International Centre for Theoretical Physics, I-34100, Trieste, Italy

d Departamento de Física, Centro de Investigación y de Estudios Avanzados del I.P.N., Apdo. Post. 14-740, 07000, México, D.F., Mexico

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Abstract

The onset of inflation in hybrid models require fine tuning in the initial conditions. The inflaton field should have an initial value close to the Planck scale $M_P$, whereas the auxiliary “orthogonal” field must be close to zero with an extreme accuracy. This problem can be alleviated if the orthogonal field decays fast into some states not coupled to the inflaton. Natural candidates for such states can be the right-handed neutrinos. We show that a non-trivial evolution of the classic sneutrino fields after inflation offers an interesting mechanism for generating a correct amount of lepton asymmetry, which being reprocessed by sphalerons can explain the observed baryon asymmetry of the Universe. Our scenario implies interesting bounds for the neutrino masses in the context of seesaw mechanism.

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1. Introduction

The hybrid inflation model [1] provides an attractive possibility for solving a range of cosmological problems [2]. It has an unique dynamical feature due to interplay between two scalar fields, inflaton $\sigma$ and an auxiliary field, so called orthogonal scalar $\phi$ which is actually responsible for providing the potential energy throughout the inflationary phase. Quite remarkably, a proper form of the inflaton potential can be naturally realized in the context of supersymmetric theories [3,4]. The simplest supersymmetric model for hybrid inflation is based on the following superpotential including the inflaton superfield $S$ and the auxiliary field $\Phi$ [3]:

$$W = -\Lambda^2 S + \lambda S \Phi^2,$$

where $\Lambda$ and $\lambda$ are the model parameters. The latter are constrained by COBE normalization for the density perturbations as $\Lambda \approx 6.5 \times 10^{16} \epsilon^{1/4}$ GeV [2], or equivalently

$$\Lambda \approx 1.3 \times 10^{15} |\eta| \lambda^{-1/2} \text{ GeV},$$

where $\epsilon, \eta \ll 1$ are the slow roll parameters.

The scalar components of these superfields have a potential

$$V(\sigma, \phi) = \Lambda^4 - \lambda \Lambda^2 \phi^2 + \frac{\lambda^2}{4} \phi^4 + \lambda^2 \sigma^2 \phi^2,$$

which has a global (supersymmetric) minima at $\sigma = 0$ and $\phi^2 = 2\Lambda^2/\lambda$. However, for $\phi = 0$ and $\sigma > \sigma_c = \ldots$
\( \frac{\Lambda}{\sqrt{\lambda}} \), the potential is flat along the \( \sigma \) axis with non-zero vacuum energy \( V(\sigma, 0) = \Lambda^4 \). This flat direction is lifted by radiative corrections resulting from the supersymmetry breaking and perhaps also by some supergravity corrections \([3]\) that give an appropriate slope to the inflaton potential on which \( \sigma \) can slowly roll down and produce inflation. During inflation the orthogonal field is held in a false vacuum \( \phi = 0 \), while the inflaton must evolve from an initial value close to the Planck scale \( M_P = 1.2 \times 10^{19} \text{ GeV} \). The inflationary phase ends when the inflaton field approaches the critical value \( \sigma_c \), after which both fields \( \sigma \) and \( \phi \) begin to oscillate near their global minima and reheat the universe.

However, as any inflationary paradigm, hybrid inflation models also suffer from the initial condition problems. A generic problem concerns the difficulty to inflate an arbitrary space–time patch. In order to trigger inflation, the inflaton field should be extremely homogeneous in an initial patch of the Universe at scales larger than the Hubble radius of this epoch \([5]\). This becomes a challenging task, since we know from the observations that the present inhomogeneity is primordial in nature and necessarily ties its origin with inflation. This obstacle can be evaded for a chaotic initial condition provided the fields take values close to the Planck scale.

In the case of hybrid inflation this generic problem is aggravated by the fact that strong fine tuning is required also for the initial values for the fields \([6,7]\). This consists in the following. At the Planck epoch, one can expect that all scalar fields, not only the inflaton, take their initial values close to the \( M_P \). However, if the orthogonal field has a value \( \phi \sim M_P \), it does not provide the right condition for the onset of inflation. The reason is simply the following. While the field \( \sigma \sim M_P \) provides large effective mass term to \( \phi \) through the last term in (3), the value \( \phi \sim M_P \) in turn would induce large mass term to \( \sigma \). In this case both the auxiliary field and inflaton would merely oscillate at their local minima without inflating the Universe. In order to trigger inflation, the initial field configuration should be settled down along the \( \sigma \)-valley, with \( \sigma \) having the value of order \( M_P \) while \( \phi \) must be close to zero with extreme accuracy, \( \phi < 10^{-5} M_P \) \([6,7]\). This is precisely the point which is not very natural; why the orthogonal field should start so close to the false minimum of its potential? In other terms, in order to inflate the initial patch of space after the Planck epoch, it should not only be strongly homogeneous at distances much larger than the corresponding Hubble radius, but it also should have energy density \( \sim \Lambda^4 \), about 15 orders of magnitude smaller than \( M_P \).

Certainly, such stringent initial conditions can be accepted on purely anthropic grounds. However, it is always appealing to obtain a more natural solution of the problem. It has been shown recently in Ref. \([7]\), that the supersymmetric hybrid model can allow a solution to this fine-tuning problem quite elegantly. The solution is simple as it assumes that all the fields take their initial value close to the Planck scale. However, the orthogonal field \( \phi \) must decay into the quanta of some extra fields which do not interact with the inflaton, and thereby do not contribute to the curvature of the inflaton potential. This can be easily achieved by adding to the superpotential (1) a term \( \kappa \Phi \Psi^2 \), where \( \Psi \) is the extra superfield. The field \( \phi \) oscillating near its false minimum, which one can consider as a matter dominated phase before the onset of inflation, decays into \( \Psi \) and settles down to zero quickly enough such that the value of the inflaton remains \( \sigma \sim M_P \) which allows inflation to begin.

The graceful exit of inflation take place when the inflaton crosses the critical point on the valley which unfolds the strong positive curvature of the false vacuum of an auxiliary field to a smooth potential with a negative curvature which allows this field to roll down to its true minimum. When this happens both the inflaton and the auxiliary field begin oscillations around their respective global minima. One of the interesting aftermath of any inflationary dynamics is that the fields which take part in inflating the Universe, produce entropy and a thermal bath during their oscillations, the process known as the reheating of the Universe.

Shortly after, or possibly during reheating another crucial event, must occur: baryogenesis. As far as any preexisting baryon asymmetry has been exponentially diluted during inflation, one has to find a valid mechanism for generating the observed baryon asymmetry.

By now there are plenty of models for baryogenesis. One of the attractive possibilities would be to produce the baryon asymmetry just during the reheating, in an out of equilibrium decay of the inflaton particles themselves. Notice, that the departure from the thermal equilibrium is one of the three basic criterion for
baryogenesis other than \( CP \) and \( B \) (or rather \( B - L \)) violation [8]. It is not easy, however, to naturally realize this possibility in the context of supersymmetric hybrid inflation.

Another interesting idea is related to leptogenesis [9], where first the non-zero lepton number \( L \) is generated in some out-of-equilibrium processes, usually assumed as a right-handed neutrino [10], and it is partially reprocessed into baryon number \( B \) via \( B + L \) violating sphaleron effects which however preserve \( B - L \) [11].

Yet another approach, known as Affleck–Dine mechanism [12], takes advantage of the flat directions in supersymmetry which carry \( B - L \) global charge. The non-zero baryon number density can be produced in decay of these modes at later times in the Universe evolution.

Baryogenesis mechanism which we propose in this paper is a blend of these three ideas discussed above, and it can be naturally realized in the context of the supersymmetric hybrid inflation model [7].

Namely, a new step which we make here is to identify the extra superfield \( \Psi \), needed for alleviating the problem of initial conditions, as a right-handed (RH) neutrino. First of all, this proposal can be motivated by the observation that the inflationary energy scale \( \Lambda \sim 10^{15} \) GeV is also of interest as the RH neutrino mass scale in the context of the seesaw mechanism [13]. Indeed, in the global minimum the field \( \phi \) receives a vacuum expectation value (VEV) and thus the coupling \( \Phi \Psi^2 \) induces the mass \( M_\phi \sim \Lambda \).

Second, in this case is the scalar component of this superfield, the RH sneutrino field \( \Phi \), carries the lepton number. The associated exact global symmetry, which is actually \( U(1)_{R-L} \), is violated by the VEV of \( \phi \). However, during the inflation the system is trapped in a false vacuum with \( \phi = 0 \), where the \( U(1)_{R-L} \) symmetry is restored and \( \tilde{\Phi} \) behaves as massless field. Along with \( \sigma \) and \( \phi \), also \( \tilde{\Phi} \) must have initial value \( \sim M_\phi \) and before inflation it oscillates around zero. However, once the field \( \phi \) decays, \( \tilde{\Phi} \) becomes essentially massless mode. Therefore, during the de Sitter phase its evolution is slow and finally at the end of inflation it has non-zero value of the order of the Hubble parameter \( H \sim \Lambda^2/M_\phi \). In postinflationary epoch this field becomes massive and starts to oscillate near the origin. As we show below, their evolution in this epoch could generate, due to dynamical breaking of the associated global \( U(1) \) symmetry, an adequate amount of \( B - L \) in the Universe, which can be converted into the baryon asymmetry via sphaleron transitions.

The Letter is organized as follows. We begin with presenting our model. In next section we discuss the evolution patterns of the scalar fields before, during, and, after inflation and the reheating of the Universe. Then we calculate the lepton asymmetry generated by the RH sneutrino field and discuss seesaw phenomenology for neutrino masses. Finally, we conclude with a brief discussion of our findings.

2. The model

From the point of view of the particle physics, our model is nothing but the MSSM including the standard fermion superfields: leptons \( l = (\nu, e) \) and quarks (which we do not write explicitly), two Higgs doublets \( \varphi_{1,2} \), and in addition an extra superfield \( \Psi \) to be identified with the right-handed neutrino, which is a gauge singlet of the standard model. For simplicity, let us consider only one fermion generation.

The charged lepton gets mass from the Yukawa term \( hle \varphi_1 \), while the similar coupling \( g\Psi \varphi_2 \) induces the Dirac mass for the neutrino component. For implementing the hybrid inflation scenario, we consider the superpotential of the following form [7]:

\[
W = -\Lambda^2 S + \lambda S \Phi^2 + \kappa \Phi \Psi^2, \tag{4}
\]

which is the simplest modification of the original one (1). It is essential the term \( \kappa \Phi \Psi^2 \) which communicates the standard particle sector to the fields producing inflation. The above superpotential as well as the Yukawa terms respect the global \( R \)-symmetry with the fermion superfields \( \Psi, l, e \), etc., carrying the \( R \) charge 1/2 and the Higgs ones \( S \) and \( (\Phi, \varphi_{1,2}) \) carrying charges 1 and 0, respectively. Notice that this \( R \) symmetry forbids the \( R \)-violating terms \( l\varphi_2, ll\varphi_2 \), etc., in other words, the theory has an automatic matter parity under which all the fermion superfields change sign while the Higgs ones remain invariant. Interestingly, \( R \) symmetry forbids also the supersymmetric mass term \( \mu \varphi_1 \varphi_2 \) which seems to be a good starting point for solving the hierarchy problem. Needless to say, we assume that the Higgsino mass can emerge in some ways
in consequence of the supersymmetry breaking. Moreover, in global supersymmetry limit, the only mass scale in the theory is $\Lambda$ in (4).

Let us take the scalar field components in a form:

$$S = \frac{\sigma}{\sqrt{2}}, \quad \Phi = \frac{\phi}{\sqrt{2}}, \quad \bar{\psi} = \frac{\psi_1 + i \psi_2}{\sqrt{2}},$$

(5)

where $R$-symmetry transformation has been used to make $S$ field real, and the imaginary part of $\Phi$ has been put to zero for simplicity. (Sometimes it is convenient to present the last field in terms of its modulus and phase: $\bar{\psi} = (\rho/\sqrt{2}) \exp(i\delta)$.) Their potential reads:

$$V = V(\sigma, \phi) + \kappa^2 \phi^2 (\psi_1^2 + \psi_2^2) + \frac{\kappa^2}{4} (\psi_1^2 + \psi_2^2)^2$$

$$+ \kappa \lambda \sigma \phi (\psi_1^2 - \psi_2^2),$$

(6)

where $V(\sigma, \phi)$ is given by (3). It has a global minimum with a non-zero VEV ($\langle \Phi \rangle$):

$$\langle \sigma \rangle = 0, \quad \langle \psi_{1,2} \rangle = 0,$$

$$\langle \phi \rangle = \phi_0 = \sqrt{2} \lambda^{-1/2} \Lambda,$$

(7)

At this minimum the superfields $S$ and $\Phi' = \Phi - \langle \Phi \rangle$ have equal masses

$$M_\sigma = M_\phi = 2\lambda^{1/2} \Lambda \simeq \eta \cdot 2.5 \times 10^{15} \text{ GeV},$$

(8)

while the last coupling in (4) induces the mass of $\bar{\psi}$:

$$M_\bar{\psi} = \frac{2\kappa \Lambda}{\lambda^{1/2}} \simeq \frac{\kappa \eta}{\lambda} \cdot 2.5 \times 10^{15} \text{ GeV},$$

(9)

where (2) has been used for numerical estimations. Interestingly, this is just the right range for the RH neutrino mass. Then, by means of the seesaw mechanism, the ordinary neutrino gets small Majorana mass

$$m_\nu = \frac{g^2 \langle \psi_2 \rangle^2}{M_\bar{\psi}} \simeq \frac{g^2 \lambda}{\kappa \eta} \cdot 1.2 \times 10^{-2} \text{ eV},$$

(10)

where for numerical estimation we have taken $\langle \psi_2 \rangle \simeq 170$ GeV (i.e., not very small $\tan \beta$).

The superpotential (4) provides the same pattern for the inflation as the original one (1). For the field values $\sigma > \sigma_0 = \lambda^{-1/2} \Lambda$, the potential (6) has a flat direction along the $\sigma$ axis, in which the orthogonal field is trapped in false vacuum with the energy $V(\sigma, 0) = \Lambda^4$. This flat direction is lifted by radiative corrections resulting from the supersymmetry breaking by the non-zero vacuum energy, and perhaps also by some other supergravity corrections [3]. The precise form of these corrections is not important, and we can simply assume that they result in an effective $\sigma$ dependent potential, e.g., in the form of mass term, $\sim \bar{m}^2 \sigma^2$, $m \ll \Lambda$, that gives an appropriate slope to the inflaton potential on which $\sigma$ can slowly roll down and produce inflation. For achieving this, the parameters

$$\epsilon = \frac{M_\phi^2}{16\pi} \left( \frac{V'}{V} \right)^2 \simeq \frac{\lambda^2 M_\phi^2 \sigma^2}{\pi \Lambda^4} \left( \frac{\bar{m}}{M_\sigma} \right)^4,$$

$$\eta = \frac{M_\bar{\psi}^2}{8\pi} \left( \frac{V''}{V} \right) \simeq \frac{\lambda M_\phi^2}{2\pi \Lambda^2} \left( \frac{\bar{m}}{M_\sigma} \right)^2,$$

(11)

have to be much smaller than one. For producing inflation, the orthogonal field should be settled in a false vacuum $\phi = 0$ while the inflaton evolves from an initial value $\sigma \sim M_\phi$. The inflationary phase ends up when the inflaton approaches the critical value $\sigma_c$, after which both fields $\sigma$ and $\phi$ begin to oscillate near their global minima and reheat the universe.

Let us now turn to the role of the superfield $\bar{\psi}$. In the global minimum (7), its fermion component, the RH neutrino, gets large Majorana mass and thus violates the lepton number conservation. However, the Lagrangian of the scalar component $\bar{\Psi}$, the RH sneutrino, maintains the global $U(1)$ symmetry. Indeed, since $\sigma = 0$, the scalar potential (6) becomes a function of the field $\bar{\Psi}$ modulus $\rho = (\psi_1^2 + \psi_2^2)^{1/2}$ and does not depend on its phase $\delta$. Needless to say, that the Yukawa coupling to lepton and Higgsino, $\bar{\Psi} l \bar{\psi}_2 + \text{h.c.}$, are conserving the lepton number.

On the other hand, in the false minimum $\phi = 0$, the superfield $\bar{\psi}$ is massless and completely decoupled from the inflaton $\sigma$. Now the scalar potential of $\bar{\Psi}$ contains only the last term $\propto \kappa^2 \rho^4$ in (6). Thus, the global $U(1)$ symmetry associated with the lepton number is restored.

Therefore, the $U(1)$ symmetry in the potential of RH sneutrino fields $\bar{\Psi}$ can be violated only by the last term in (6), when that both $\phi$ and $\sigma$ have non-zero values. This can occur only during the epoch when the background fields $\phi$, $\sigma$ oscillate near their global minima. As we show below, at this phase the non-zero lepton number can be generated in the classical motion of the RH sneutrino fields, which will be transferred to the standard leptons via their decay $\bar{\Psi} \rightarrow l \bar{\psi}_2$. 

3. Dynamics of the fields

In this section we shall describe the evolution of the fields in our model. Starting our consideration from an earliest time when the Planck era ends and classical general relativity becomes applicable, we assume that at this moment all classical fields have initial values $\sim M_P$ and hence the energy density of the Universe is $\sim M_P^4$, quite a generic situation. We study the scalar field dynamics before the inflation onset, during exponential expansion, their postinflationary oscillations and, finally, we discuss the reheating of the Universe. This will provide us all necessary tools to calculate the lepton asymmetry.

3.1. Pre-inflation

The detailed analysis of the scalar field dynamics before inflation in the model with the superpotential (4) has been performed in Ref. [7]. Here we briefly recall its main features.

All scalar fields $\sigma$, $\phi$ and $\psi_{1,2}$ have initial values close to the Planck scale, and with the Hubble expansion they evolve down due to interaction terms in (6). Therefore, the cosmological energy density, initially $\sim M_P^4$, decreases with the Hubble expansion and its evolution settles quickly in a pattern when the energy of the system is dominated by the regular oscillations of $\phi$ around zero. From this moment the Universe expands as in a matter dominated era.

Along the flat direction the effective mass for $\phi$ field is quite large, $M_{\phi,\text{eff}}^2 = 2\lambda^2(\sigma^2 - \sigma_c^2)$, and it vanishes only when $\sigma = \sigma_c$. On the other hand, once $\phi$ has a large amplitude, it induces the large curvature for the inflaton field $\sigma$ and the latter starts to quickly roll down to the origin. Therefore, in order to produce inflation, the amplitude of $\phi$ must settle to zero before $\sigma$ will manage to substantially decrease its initial value $\sim M_P$. The friction provided by the Universe expansion itself cannot really help, since the Hubble parameter in this epoch is actually smaller than $M_{\phi,\text{eff}}$ and thus the field $\phi$ suffers many oscillations during the Hubble time.

However, the problem can be with help of the coupling $\kappa \Phi \psi^2$, which permits the oscillating field $\phi$ to decay into $\Psi$ particles. The decay rate $\Gamma(\phi \rightarrow \Psi \Psi) = (\kappa^2 / 8\pi)M_{\phi,\text{eff}}$ can be large as far as the coupling constant is reasonably large, $\kappa > 0.1$ or so, and $\sigma \sim M_P$. In addition, as soon as $\phi$ becomes zero, $\sigma$ stops to feel the classical fields $\psi_{1,2}$ as well as their quanta produced in the decay, large curvature of the inflaton potential disappears and its value freezes at values $\sigma \sim M_P$. By the time $t \sim 1 / \Gamma$ the cosmological energy density becomes dominated by the relativistic $\Psi$ particles, and it decreases fast with the universe expansion until it becomes dominated by the false vacuum energy. Since this moment Universe starts to expand exponentially while the inflaton proceeds its slow roll due to small curvature term.

The decay of $\phi \rightarrow 2\Psi$ can be understood as a process of preinflationary reheating after which the system from a generic initial state, with all fields having magnitudes $\sim M_P$, settles in false vacuum. Therefore, it provides a simple mechanism to obtain the correct initial condition for the inflation to take place. Once inflation commences it dilutes the produced $\Psi$ quanta.

Let us turn now to the evolution of the classical sneutrino field $\tilde{\psi}$ (5). As far as we did not introduce any mass term for the RH neutrino, in the superpotential (4), the fields $\psi_{1,2}$ are intrinsically massless. Nevertheless, before the decay of $\phi$ oscillations, they have field dependent effective mass terms $m_{\phi,\text{eff}}^2 = \kappa^2 \phi^2 \pm \kappa \lambda \sigma \phi$ and therefore during the matter dominated phase prior to inflationary phase they simply roll down from the initial values order $M_P$. However, when $\phi$ settles to zero and inflation begins, these fields become massless and their oscillations are damped. By this moment these fields still have reasonably big values. However, they have further non-trivial evolution during and after inflation, which we shall discuss in subsequent sections.

3.2. Inflation

As was told in the above, the inflation proceeds when $\phi$ is set to zero while $\sigma$ still has a large value $\sim M_P$. On the other side, $\sigma$ field has a small effective mass term $\sim m^2\sigma^2$ which allows it to roll down the potential. The main contribution to the energy density comes from the false vacuum, $V(\sigma, 0) \simeq \Lambda^4$, and the Universe expands exponentially up to the time $t$ when the inflaton field reaches the critical value $\sigma_c$, where on the Universe exits gracefully from almost de-Sitter expansion during which the cosmological scale has...
grown by a factor $\exp(\mathcal{H} t)$. The Hubble parameter during the inflation is given by

$$ H \simeq \sqrt{\frac{8\pi}{3} \frac{\Lambda^2}{M_p}} \simeq \frac{\eta^2}{\lambda} \cdot 4 \times 10^{11} \text{GeV}. $$

(12)

As soon as $\phi$ settles to zero and inflation begins, $\psi_{1,2}$ become massless and the curvature of their potential is induced only by quartic self-couplings in (6). These fields are completely decoupled from the inflaton and thus they evolve almost independently during the inflationary era, with the following equations of motion:

$$ \ddot{\psi}_{1,2} + 3H \dot{\psi}_{1,2} = -\kappa^2 (\psi_1^2 + \psi_2^2) \psi_{1,2}, $$

(13)

or, in terms of the modulus $\rho$, $\dot{\rho} + 3H \dot{\rho} = -\kappa^2 \rho^3$. As for the phase $\delta$, it essentially becomes a flat mode as far as $\phi = 0$ and the lepton number conservation is restored.

The classical field $\rho$ keeps rolling down until it approaches values $\sim H$. After this its dynamics is almost frozen and the rest of evolution is determined by the equation $3H \dot{\rho} \approx -\kappa^2 \rho^3$. From here one immediately obtains that by the time when the slow roll ends up, the field value will be $\rho \approx \kappa^{-1} (3H/t)^{1/2}$. This can be rewritten as

$$ \rho = \frac{\sqrt{C}}{\kappa} H, $$

(14)

where $C = 3/N_c$, with $N_c = Ht$ being the total number of e-foldings. Notice also, that in hybrid inflation models, unlike chaotic models, the number of e-foldings can not be arbitrarily large and for our purpose we assume that maximum number of e-folding can at most be $N_c \approx 100$.

So far in the evolution of $\Psi$ fields we have not taken into consideration any kind of supergravity correction to their masses. In particular, for a generic Kähler potential the fields $\Psi$ could get the mass term $-C H^2(\psi_1^2 + \psi_2^2)$, $C$ being order 1 coefficient. The sign of $C$ is model dependent and it can not be determined correctly. If the correction is positive, then during inflation $\rho$ evolves as $\propto e^{-3\mathcal{H} t}$ and it will essentially vanish at the inflation exit. However, if the mass correction turns out to be negative, then $\Psi$ fields will have a false minimum with a non-zero VEV which breaks the lepton number: $(\rho)^2 = C H^2/\kappa^2$. Therefore, during inflation these fields will fast evolve down until being trapped in the false minimum, and remain stuck to it until the effective mass of $\Psi$ field overtakes the expansion parameter $H$. This happens very soon after the end of inflation, because $M_{\Psi}^2 \propto \Lambda^2 \gg H^2$. Therefore, by the end of the inflation era $\rho$ has a non-vanishing magnitude which can be given still by (14), but this time $C$ being some unknown order 1 coefficient.

### 3.3. Post-inflation

The evolution of $\sigma$ and $\phi$ after the inflation exit has been studied in Refs. [14,16]. After the phase transition the fields oscillate with more or less similar amplitude near their global minima (7), the maximum amplitude attained by $\sigma$ field is $\sigma_c = \phi_0/\sqrt{2}$, while $\phi$ takes at most $\phi_0$. The initial conditions for the oscillations are fixed by the inflationary dynamics and are given by [14,16]:

$$ \sigma_i = \sigma_c \pm \frac{H_i}{\sqrt{2}}, \quad \dot{\sigma}_i = -\frac{1}{3H_i} \frac{\partial V}{\partial \sigma}, $$

$$ \phi_i = \frac{H_i}{2\pi}, \quad \dot{\phi}_i = -\frac{1}{3H_i} \frac{\partial V}{\partial \phi}, $$

(15)

where the initial velocities are calculated at $\phi_i$ and $\sigma_i$, and $H_i$ is the Hubble parameter at the end of inflation.

The oscillations are essentially anharmonic in nature. However, near the global minimum, there exists a particular solution which satisfies a straight line trajectory in $\phi$–$\sigma$ plane [14,16];

$$ \phi = \sqrt{2}(\sigma_c - \sigma). $$

(16)

Near the bottom of the potential the oscillations are harmonic and their frequency is governed by the mass of the fields at their minima, see (8). It is possible to give an approximate analytical solution

$$ \frac{\dot{\phi}(t)}{\phi_0} \approx 1 + \frac{A(t)}{3} \cos(m_\phi t), $$

(17)

where $A(t) \sim 1/t$ is slowly time-varying amplitude of the oscillations, which depends crucially upon the ratio $H_i/M_\phi$. The smaller the ratio is, the larger is the number of oscillations of the fields before they feel the Hubble expansion. For smaller inflationary scales, such as $A \ll 10^{-3} M_p$, the ratio $H_i/M_\phi$ is small, $H_i/M_\phi \simeq 1.3 A/\sqrt{3} M_p \ll 1$. This is an interesting characteristics of the supersymmetric hybrid inflation model, which tells us that there are many oscillations of the background fields with an almost constant
amplitude. The effect of expansion is felt after many oscillations, and this has been verified numerically in Refs. [14–16].

Let us consider now post-inflationary dynamics of the classical fields $\Psi$. Although they have been essentially the massless fields during the inflation, after inflation they become massive, with mass $M_\Psi \sim \Lambda$. In other words, $\psi_{1,2}$ fields get quadratic terms in the potential and they start to oscillate near the origin with the frequencies governed by mass $M_\Psi$ (9). Depending on the situation, the initial amplitudes for these oscillations, see (14). For keeping more generality, let us parameterize the magnitudes of these amplitudes as

$$\psi_{1,2}^i = \rho_i \cos \delta_i (\sin \delta_i), \quad \rho_i^2 = \frac{C}{k^2} H, \quad (18)$$

where we assume that the phase $\delta$ is order one, and $C = 3/N_e$ or $C \sim 1$, depending on the situation whether at inflation stage these fields had the supergravity induced order $H^2$ mass terms or not. Therefore, the classical fields $\psi_{1,2}$ have the initial amplitudes $\sim H$, while $\sigma$ and $\phi$ fields have much larger amplitudes $\sim \Lambda/\sqrt{\lambda}$. The initial energy density $V \simeq \Lambda^4$ is dominated by oscillations of $\sigma$ and $\phi$, while the contribution of $\Psi$ is negligible. In this way, one can safely neglect backreaction of $\Psi$ fields on the oscillation of the classical background fields $\sigma$ and $\phi$.

The post-inflationary evolution of classical $\Psi$ fields is quite interesting. The oscillation frequency of these fields is governed by the value of $\phi$, and also the evolution of $\sigma$ leads to an additional contribution to their equation motion. On the other hand, once they become massive and can decay into leptons and higgsinos, the decay rate (22) also contributes the friction term in their equation of motion, which now read

$$\ddot{\psi}_{1,2} + (3H + \Gamma) \dot{\psi}_{1,2} = -k^2 (2\dot{\phi}^2 + \psi_{1,2}^2 + \psi_{1,2}^2) \psi_{1,2} \mp 2k\lambda \sigma \phi \psi_{1,2}. \quad (19)$$

Notice, that the last terms in the above equations come with an opposite sign and so the evolution of $\psi_1$ is different from $\psi_2$. In other words, in the background of the $\sigma, \phi$ fields, $\psi_{1,2}$ get the $U(1)$ invariant effective mass term $M_\psi^2(\sigma \phi)\Psi^* \Psi = 2k^2\dot{\phi}^2(\psi_{1,2}^2 + \psi_{1,2}^2)$, as well as the $U(1)$ violating one $M_\psi^2(\sigma \phi)(\dot{\psi}_{1,2}^2 + \dot{\psi}_{1,2}^2) = 2\lambda k \sigma \phi (\psi_{1,2}^2 - \psi_{1,2}^2)$.

As we have already remarked, $\phi, \sigma$ fields make many oscillations in one Hubble time, which allows us to consider the average effect of these fields upon $\psi_1, \psi_2$ fields. In other words, one can replace $\phi^2$ and $\sigma \phi$ by their mean values within a period of one Hubble time, for which from (16) and (17) we obtain:

$$\langle \phi^2 \rangle \simeq \frac{2\Lambda^2}{\lambda}, \quad \langle \sigma \phi \rangle \simeq \frac{\sqrt{2} \Lambda^2}{18\lambda} A^2(t). \quad (20)$$

where $A(t) \propto 1/t^2$. Substituting these averages in the (19), we see that during postinflationary oscillations $\psi_1$ and $\psi_2$ have different dynamical mass terms:

$$M_{\psi_1}^2(t) = M_{\psi_1}^2 \pm M^2 = M_\phi^2 \left(1 \pm \frac{A(t)^2}{25}\right). \quad (21)$$

This dynamical mass splitting causes their helical motion in the background of $\sigma$ and $\phi$ and produces the lepton asymmetry in our model.

The evolution $\Psi$ fields crucially depends whether the friction term in (19) is dominated by $H$ or $\Gamma$. In the former case, the amplitude $\psi$ decreases with time as $\propto 1/t$, while in the latter case as $\exp(-\Gamma t/2)$. For the width of the decay $\Psi \rightarrow l\bar{\nu}_2(l\psi_2)$ we have:

$$\Gamma_\Psi = \frac{g^2}{8\pi} M_\Psi \simeq \left(\frac{\kappa \eta}{\lambda}\right)^2 \left(\frac{m_\nu}{0.1 \text{eV}}\right) \times 10^{15} \text{GeV}. \quad (22)$$

### 3.4. Reheating

The end of inflation also marks an era of entropy creation in the Universe. The energy stored in oscillations of classical fields $\phi$ and $\sigma$ eventually decays into relativistic particle species. In our case reheating could occur via two possible channels. Either $\sigma$ and/or $\phi$ have some dominant channel to decay into standard particles, or, $\phi$ could decay into right-handed neutrino $\Psi$ quanta, which would subsequently decay into leptons and Higgs to produce a final thermal bath. The choice of the dominant channel shall depend on the couplings in the theory which in our case are constrained to guarantee that the reheat temperature of the Universe is not too large. In particular, in supersymmetric theories there is an upper bound on the reheating temperature $T_r \leq 10^9$ GeV or so [17]. If this condition is not met, then the gravitinos are effectively produced via scattering processes in thermal bath, and their late decay products can genuinely threaten the
The reheating temperature of the Universe era can be estimated as

$$T_r \approx 0.5 \, g_*^{-1/4} \sqrt{\Gamma M_P} \simeq 0.1 \sqrt{\Gamma M_P}, \quad (23)$$

where $\Gamma$ is the decay rates of the fields $\sigma$ or $\phi$, and $g_*$ is an effective number of the particle degrees of freedom in the thermal bath. (In the supersymmetric standard model which we consider here, $g_* \sim 200$).

The following remark is in order. Apart from the perturbative decay of the oscillating fields, the coherent oscillations of the inflaton could also lead to a non-thermal resonant production of particles [18], and in particular of gravitinos with helicity 3/2 [19], and helicity 1/2 [20]. This would produce very stringent upper bound on the reheating temperature. However, it has been realised later on that non-thermal production of gravitinos during inflation occurs via the goldstino mode, a helicity $\pm 1/2$ component which can be recognised as inflatino. It eventually decays along with the inflaton to reheat the Universe, and thus non-perturbative production of gravitinos should not pose serious problems for nucleosynthesis [21].

One of the most important criteria for a resonant particle production is the coherent oscillations of the background fields, which is fulfilled in supersymmetric hybrid inflationary model. It has been noticed that the particle creation is quite efficient in supersymmetric hybrid models, see Ref. [14]. In particular, quanta of the heavy RH neutrinos produced at the preheating stage by non-perturbative decay of inflaton oscillations, in their subsequent decays could produce the lepton asymmetry, provided that these decays are CP-violating [22].

Another interesting feature of hybrid model which we must mention here is the possibility to realise tachyonic preheating. This is because near the critical point mass squared for $\phi$ field flips its sign and becomes a tachyonic mode. This violates the adiabatic vacuum condition and leads to explosive production of particles [23]. In such a case, it is also important to consider the backreaction of the quanta on the classical fields, which actually leads to destabilising the system and the zero-mode trajectories [16]. Eventhough, tachyonic preheating might work, it is still not very clear why there should be a resonant particle production once the backreaction of the quanta is properly taken into account.

Whatevver be the case, we shall not worry too much upon the preheating aspects. All we assume here is that the decay of $\phi$ and $\sigma$ is responsible for reheating the Universe. We also neglect the finite temperature effects, which could emerge if the relativistic particle species produced by the inflaton decay thermalize too early, before the inflaton decay ends up, and thus can give thermal corrections to the potential of the fields such as $\Phi$ in our case. The finite temperature mass corrections can be avoided once thermalization of the Universe is delayed until the last stages of reheating.

Coming back to our situation, we have to control that the already existing couplings in the model will not produce too large reheat temperature. The dominant channel for the decay of $\Phi$ field is governed by the last term in the superpotential (4). If $M_\phi > M_\Psi$, then $\Phi$ decays into $\Psi$ quanta which then produce the standard particles via the coupling $g_\Psi \bar{l} \varphi$. In this case we have $T_r = (\kappa^2 / 8 \pi) M_\phi$, and thus $T_r \sim 0.1 \sqrt{T_0 M_P} \simeq 3 \kappa \eta \times 10^{15}$ GeV, where we have used (8). Taking now into account the fact that in our approach the constant $\kappa$ should be rather large, since otherwise the decay of the orthogonal field $\phi$ would not be effective in the preinflationary phase, and neither $\eta$ can be be very small, we obtain too a big reheat temperature.

However, if $M_\Psi > M_\phi$, i.e., $\kappa > \lambda$, situation is more interesting and the problem can be solved without any extra assumptions. Now the decay channel $\Phi \to \Psi \bar{\Psi}$ is kinematically forbidden, and the lowest order relevant operator for the decay of $\phi$ into lighter particles is the $D = 6$ one (e.g., $g_\phi^2 / M_\phi^2) \Phi \bar{l} \varphi$. Such term in the superpotential is induced by the Yukawa couplings $g \bar{\Psi} l \varphi_2$. Integrating out the heavy superfields $\Psi$, the decay width of the four-body decay $\Gamma (\phi \to ll\varphi) \sim \kappa^2 g^4 M_\phi^2 / M_\phi^2$ is suppressed by a phase space factor $P \sim 10^{-5}$. Finally, by taking into account (8), (9) and (10), the reheating temperature comes naturally as

$$T_r \simeq \lambda \eta^{3/2} \left( \frac{m_\varphi}{0.1 \text{ eV}} \right) \times 10^{14} \text{ GeV}. \quad (24)$$

For instance, by taking both $\eta$ and $\lambda$ order $10^{-2}$ and $m_\varphi \sim 0.1$ eV, we get $T_r \sim 10^9$ GeV. This result also satisfies the thermal gravitino production bound on the
reheating temperature. In addition, as we shall see in the next section, such a reheat temperature can easily produce the right amount of the baryon asymmetry.

4. Lepton asymmetry

As we have mentioned earlier, the RH sneutrino field $\tilde{\Psi}$ carries the global $U(1)$ charge, and the associated quantum number $B - L$ (rather than $L$) is violated via the last term in the potential (6). This term is effective during the post-inflationary oscillations and it gives rise to a helical motion of $\psi_{1,2}$ in the background of $\sigma$ and $\phi$ fields, thereby generating the lepton number in the postinflationary universe.

The $B - L$ charge density generated during the time evolution of the RH sneutrino field is a zero component of the global $U(1)$ current:

$$n_{B-L} = \frac{i}{2}(\dot{\psi}_1^*\tilde{\psi} - \dot{\psi}_2^*\tilde{\bar{\psi}}) = \frac{1}{2}(\psi_1\dot{\psi}_2 - \dot{\psi}_1\psi_2).$$

(25)

Then from (19) we immediately obtain the equation which describes its evolution:

$$\dot{n}_{B-L} + 3Hn_{B-L} = 2\kappa\lambda\sigma(t)\phi(t)\psi_1(t)\psi_2(t),$$

(26)

where the RH side acts as a source term which generates a net $B - L$ asymmetry through a non-trivial motion of $\psi_1$ and $\psi_2$ fields.

Let us integrate this equation from the time moment $t_i$, which corresponds to the end of inflation, up to a finite time interval $t$:

$$n_{B-L} = \frac{2\kappa\lambda}{R_i} \int_{t_i}^{t} dt' R_i^{\frac{1}{2}}(\sigma(t')) \rho^2(t') \sin 2\delta_{\psi},$$

(27)

where $R_i$ is scale factor and we have substituted the mean value $\langle \sigma(t) \phi \rangle_t$ given in (20). The $CP$-phase $\delta_\psi$ changes slowly, since for $\lambda < \kappa$ oscillations of the fields $\psi_1$ and $\psi_2$ have about the same oscillation frequency (see (21)). In terms of the field quanta, $n_{B-L}$ is nothing but a number density difference between the RH sneutrino $\tilde{\Psi}$ and anti-sneutrino $\tilde{\bar{\psi}}^*$ states. It will be transmitted to the standard particle system via the decay of the RH sneutrinos into the ordinary leptons and Higgsinos (or sleptons and Higgses). Therefore, even if the decay rates $\tilde{\Psi} \rightarrow l\bar{\nu}_2$ and $\tilde{\bar{\psi}} \rightarrow \bar{l}_2\bar{\nu}_2$, are exactly the same (no $CP$-violation in decays), we produce the different amount of $l$ and $\bar{l}$.

Once the $B - L$ is non-zero, the net baryon number is induced via sphaleron effects which violate $B + L$ but conserve $B - L$ [11]. The sphalerons are active in a temperature range from about $10^{12}$ GeV down to 100 GeV. In the context of the supersymmetric model the relation between the $B$ and $B - L$ is given by $B = -0.35(B - L)$ [24]. Therefore, for obtaining the observed baryon to entropy density ratio in the range $B = n_B/s = (0.3 - 1) \times 10^{-10}$, as it is restricted by the primordial nucleosynthesis bounds, we need $B - L \sim O(10^{-10})$.

Now we are in grade to calculate the $B - L$ number to entropy density ratio $B - L = n_{B-L}/s$ in the Universe. Assuming that the entropy is generated at the postinflationary reheating and there is no more entropy injection at later times, we calculate from (27) the value of $n_{B-L}$ produced by the reheating time $t = t_r \approx 0.3 g_s^{1/2}M_P/T_r^2$, and compare it to the entropy density at this time, $s = (2\pi^2/45)g_sT_r^3$.

The Universe dominated by the field oscillations expands as in a matter dominated era, and so the scale factor changes as $R \propto t^{2/3}$. On the other hand, $\langle \sigma(t) \phi(t) \rangle_t \propto 1/t^2$ while the $CP$-phase $\delta_\psi$ changes slowly, since for $\lambda < \kappa$ oscillations of the fields $\psi_1$ and $\psi_2$ have almost the same oscillation frequency. Therefore, the integrand in (27) goes as $\rho^2(t)$, and the final answer depends whether the decay rate (22) is larger or smaller as compared to the Hubble parameter. Let us consider first the case $\Gamma < H$. As we discussed in previous section, then we have $\rho^2 \propto 1/t^2$. This suggests that the maximum contribution to integral (27) comes at the initial times $\sim 1/H$, and hence we obtain:

$$B - L \simeq \frac{\kappa \rho_1^2}{H^2 M_P} \frac{T_i}{T_r} \sin 2\delta_{\text{eff}},$$

(28)

where $\delta_{\text{eff}} \sim 1$ is an effective $CP$-violating phase which we assume to be $O(1)$, and $\rho_1$ is given by (18).

However, if the friction term in (19) is dominated by the decay width, i.e., $\Gamma > 3H$, then the RH sneutrino field behaves like $\rho^2 \propto \exp(-\Gamma t)$ and, as a result, the magnitude of $B - L$ will be reduced by a factor $x = 3H/\Gamma$ which can be obtained from (12) and (22). Therefore, the final result for the lepton asymmetry reads:

$$B - L \simeq X \left( \frac{T_i}{10^9 \text{GeV}} \right) \times 10^{-10},$$

(29)
where the numerical factor is \( X = \min[1, x | \mathcal{K}|^{-1}] \), with \( C \) being \( \sim 1 \) coefficient if the field \( \Psi \) gets supergravity induced mass term during inflation, or \( C = 3/N_c \) otherwise.

For \( X \sim 1 - 10 \), the result (29) implies correct magnitude of the baryon asymmetry if \( T_e \sim 10^9 \) GeV or so. This range for the reheating temperature can be natural in the context of our model, provided that the factor \( \lambda \eta^{3/2} \) in (reftr) is \( O(10^{-5}) \). The larger \( T_e \) would contradict to the non-thermal gravitino production limit. On the other hand, if \( x > 1 \), then the coefficient \( X \) order 1 or 10 can easily emerge if \( \kappa \sim 10^{-1} \) and \( C \sim 1 \) or 1/20, the latter value attained to the case of 60 e-fold inflation.

5. Discussion

Up to now we have considered only one fermion generation. Let us incorporate now all three generations and discuss what happens in this case. In other words, we introduce 3 lepton species \( l_a \), \( e_a \) and \( \Psi_a \), \( a = 1, 2, 3 \) is a generation index.

Now the last term in the superpotential (4) becomes \( \kappa_{a} \Phi \Psi^2 \) (we have taken this couplings diagonal, i.e., we work directly in a basis of the RH neutrino mass eigenstates). Let us assume for simplicity, that the neutrino Dirac terms are also diagonal in this basis: \( g_{a} l_{a} \Psi_{a} \Phi \). In fact, there are the fermion mass models of this type (see, e.g., Ref. [25]), in which the neutrino mass matrix is diagonal and non-zero mixing angles in lepton sector emerge exclusively from the charged lepton mass matrix.

Let us assume that all three constants \( \kappa_{1,2,3} \) are enough large, namely \( \kappa_{a} > \lambda \), in order to evade the excessively large reheating temperature of the Universe. In this case, the hierarchy of neutrino masses goes as \( m_{\nu a} \equiv m_{\nu a} \propto g_{a}^2 / \kappa_{a} \), and it should emerge from the hierarchy of the constants \( g_{1,2,3} \).

The evolution pattern of the classic fields can be extended for the case of three RH sneutrinos in a straight forward manner. In particular, we see from (24) that the reheating temperature \( T_e \) is essentially determined by the largest neutrino mass. Recalling also that the atmospheric neutrino oscillations point to the neutrino mass in the range \( m_{3} \sim 0.1 \) eV, we see that \( T_e \sim 10^9 \) GeV can be obtained in our model provided that \( \eta \sim 10^{-2} \) and \( \lambda \sim 10^{-2} \), a quite natural parameter range in the hybrid model. However, the reheating temperature much smaller than this estimate is not very appealing since it would need unnaturally small value of \( \lambda \).

On the other hand, the amount of produced \( B - L \) crucially depends on the coefficients \( x_{a} = 3H/\Gamma_{a} \). In order to avoid too strong suppression of the result (29), at least one of the factors \( x_{a} \) should be larger than 1 or so. Interestingly, this suppression factor is inverse proportional to neutrino mass—indeed, from (12) and (22) we see:

\[
x_{a} = \frac{\lambda}{\kappa_{a}^2} \left( \frac{10^{-4} \text{ eV}}{m_{a}} \right),
\]

and thus the the largest contribution to \( B - L \) is given by the lightest neutrino mass, presumably \( \nu_{1} \). Thus, the condition \( x_{1} > 1 \) implies the upper bound on the lightest neutrino mass \( m_{1} < (\lambda/\kappa_{1}^2) \times 10^{-4} \) eV, which limit, e.g., for \( \lambda < 10^{-2} \) and \( \kappa_{1} > 0.1 \), leads to \( m_{1} < 10^{-4} \) eV. This limit can be naturally met by the the mass of the first generation neutrino, if the neutrino mass hierarchy is about the same as that of charged leptons [25].

Let us conclude by summarising some interesting features of our model, which is just the simple supersymmetric hybrid model with the superpotential (4). The couplings \( \kappa_{a} \Phi \Psi_{a}^2 \) of the auxiliary orthogonal superfield \( \Phi \) to the RH neutrino ones \( \Psi \) puts the bridge between the inflation and particle physics sectors, thus connecting the inflation scale \( \sim 10^{15} \) GeV to the RH neutrino mass scale needed in the context of seesaw mechanism. As a bonus, these terms can help in solving many problems of the inflationary cosmology and baryogenesis.

First of all, they allow the orthogonal field oscillations to decay fast enough and thus can prepare the proper initial conditions for the inflation onset starting from almost arbitrary initial field configurations with the classical field values order \( M_P \).

And second, at the epoch of postinflationary field oscillations, these terms generate the dynamical lepton number breaking for the RH sneutrino fields after the end of inflation and before the end of reheating era. During de-Sitter era these fields are intrinsically massless modes and evolve very slowly due to quartic self-couplings or because of order \( H \) mass term induced by the supergravity corrections.
In either way, at the end of the inflation these fields have non-zero values order $H$. This is a virtue of $R$-symmetry which actually forbids terms like $S\Psi^2$ or $M\Psi^2$ in the superpotential (4). After inflation they start to oscillate near origin and produce the $B-L$ asymmetry of the Universe. This happens in an elegant way because the associated $U(1)$ charge is dynamically broken in the background of the oscillating inflaton fields. After the RH neutrino decay, the produced $B-L$ number density is transferred to the standard particles, and being reprocessed by sphalerons, gives rise to the net baryon asymmetry of the Universe. In difference from the usual leptogenesis mechanisms with the RH neutrino decay [9,10,22], our mechanism does not require the presence of $CP$-violation in the lepton mixing. So, it can work in the context of predictive models [25] which do not contain these $CP$-phases but are appealing in all other respects.

In our model interesting relations emerge between the inflationary parameters, reheating temperature, $B-L$ number density and neutrino masses. The amount of the produced $B-L$ solely depends upon the reheat temperature $T_r$ with some coefficient $X$ which incorporates the coupling and can be order 1. In this case, the correct amount of the $B-L$ is obtained when $T_r \sim 10^9$ GeV, close to the upper bound from the thermal gravitino production. On the other hand, the possibility of factor $X$ to be order 1, implies the upper limit on the lightest neutrino mass (presumably $\nu_e$). In general, our model is compatible also with the neutrino mass spectrum inferred from the atmospheric and solar neutrino oscillations. One could envisage, that in the context of our observation, the hybrid models [26] designed for explaining the reheating temperature difference between the ordinary and hidden (mirror or shadow) worlds, could also generate the non-zero baryon asymmetry in both sectors.

Concluding, our model does three jobs very neatly. First, it correctly indicates the neutrino mass range by linking the inflation scale to the RH neutrino mass scale in the context of seesaw mechanism. Second, it provides dynamically the proper initial condition for the onset of inflation in hybrid model. And finally, via the classic RH sneutrino fields, it generates the proper baryon asymmetry of the universe at the epoch of post-inflationary oscillations and reheating.

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