Smooth hybrid inflation and nonthermal type II leptogenesis

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Smooth Hybrid Inflation and Non-Thermal Type II Leptogenesis

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We consider a smooth hybrid inflation scenario based on a supersymmetric $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model. The Higgs triplets involved in the model play a key role in inflation as well as in explaining the observed baryon asymmetry of the universe. We show that the baryon asymmetry can originate via non-thermal triplet leptogenesis from the decay of $SU(2)_L$ triplets, whose tiny vacuum expectation values also provide masses for the light neutrinos.

I. INTRODUCTION

There exists an attractive class of supersymmetric models in which inflation is closely linked to the supersymmetric grand unification scale [1–4]. Among these models, supersymmetric hybrid inflation (with minimal Kähler potential) predicts a scalar spectral index close to 0.985 [1], to be compared with $n_s = 0.968 \pm 0.014$ presented by WMAP7 [5]. Smooth hybrid inflation, a variant of supersymetric hybrid inflation, yields a spectral index of 0.97 if supergravity effects are ignored. However, inclusion of supergravity corrections with minimal Kähler potential leads to higher values of the spectral index even in this case [6]. It has been shown in [7, 8] that the predicted scalar spectral index in smooth hybrid inflation model is affected if the non-minimal terms in the Kähler potential are switched on, and n_s close to the WMAP prediction is easily realized. For supersymmetric hybrid inflation with soft terms, it is also possible to reduce n_s to 0.968 [9].

Inflation in these models is naturally followed by leptogenesis [10]. Type I leptogenesis from the decay of right handed neutrinos has been discussed in some details in recent papers [11], where the light neutrino masses are obtained from type I seesaw. Care has to be exercised to ensure that leptogenesis is consistent with constraints that may arise from the observed solar and atmospheric neutrino oscillations [12]. Light neutrino masses can also arise from the so-called type II seesaw mechanism [13] in which heavy scalar $SU(2)_L$ triplets acquire tiny vacuum expectation values (vevs) that can contribute to the masses of the observed neutrinos.

An interplay between type I and type II seesaw in the generation of light neutrino masses [14] is also a possibility (for example, while considering a left-right symmetric model). If the right handed neutrinos all have superheavy masses comparable to $M_{GUT} = O(10^{16}$ GeV) or close to it, the type I seesaw contribution to neutrino masses alone would be too much small to be compatible with the neutrino oscillation data. A situation similar to this is adopted in this paper where the triplet vev is the main source of light neutrino masses. It is well known that these triplet scalars can play an additional important role by producing the desired lepton asymmetry [15, 16]. They could be present in the early universe from the decay of the inflaton, and their own subsequent decay can lead to leptogenesis.

We implement this scenario (type II leptogenesis with smooth hybrid inflation) within a supersymmetric version of the well known gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [17]. (Generalizations to other (possibly larger) gauge symmetries seems quite plausible.) We restrict our attention to non-thermal leptogenesis which is quite natural within an inflationary setting. (For type II thermal leptogenesis see [18, 19]). We work here in the framework of smooth hybrid inflation [6, 20, 21] , taking into account the corrections arising from non-minimal Kähler potential. To make the scenario as technically natural as possible, we impose some additional symmetries including a $U(1)_R$ symmetry [1]. We find that the constraints from neutrino oscillations as well as leptogenesis can be satisfied with natural values of the appropriate couplings.

II. HIGGS TRIPLETS IN LEFT-RIGHT MODEL

The quark and lepton superfields have the following transformation properties under the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [17]:

$$Q = (3, 2, 1, \frac{1}{3}), \quad Q^c = (3^*, 1, 2, -\frac{1}{3}), \quad L = (1, 2, 1, -1), \quad L^c = (1, 1, 2, 1).$$

The Higgs sector consists of

$$\begin{aligned} H &= (1,2,2,0), \quad \Delta_L^a = (1,3,1,2), \quad \bar{\Delta}_L^a = (1,3,1,-2), \quad a=1,2, \\ \Delta_R &= (1,1,3,-2), \quad \bar{\Delta}_R = (1,1,3,2). \end{aligned}$$

Our primary goal, as stated earlier, is the implementation of non-thermal type II leptogenesis, and to realize it we consider two pairs of triplets Δ_L, $\bar{\Delta}_L$ (indicated by index $a = 1,2$) which, through mixing, can produce the CP violation necessary for generating an initial lepton asymmetry [30]. The model also possesses a gauge singlet superfield S which plays a vital role in inflation.

The superpotential is given by:

$$\begin{aligned} W &= S\Big[\frac{(\Delta_R\bar{\Delta}_R)^2}{M_S^2} - M_X^2\Big] + \frac{\alpha_{ab}}{M_S}\Delta_L^a\bar{\Delta}_L^b\Delta_R\bar{\Delta}_R + \frac{\gamma^a}{M_S}HH\bar{\Delta}_L^a\bar{\Delta}_R + f_1^a LL\Delta_L^a + \\ & f_2 L_c L_c \Delta_R + Y^l L L_c H + Y^q Q Q_c H, \end{aligned} \tag{1}$$

where $a, b = 1, 2$, and the $SU(2)$, generation and color indices are suppressed. M_X is a superheavy mass scale and M_S is the cutoff scale which controls the non-renormalizable terms in the superpotential. We take the matrix α_{ab} to be real and diagonal ($\alpha_{ab} = \delta_{ab}\alpha_a$) in our calculation for simplicity. The first two terms (in the square bracket) are responsible for inflation. The importance of the remaining terms will be discussed later in connection with the inflaton decay, reheating, leptogenesis and neutrino mass generation. A Z_2 symmetry along with $U(1)_R$ global symmetry is imposed in order to realize the above superpotential. The charges of all the superfields are listed in Table I. The inclusion of the Z_2 symmetry forbids terms like $\Delta_L^a\bar{\Delta}_L^b$ in the superpotential, but allows the term $\Delta_L^a\bar{\Delta}_L^b\Delta_R\bar{\Delta}_R$. This ensures that the $SU(2)_L$ triplets are lighter than the superheavy right handed neutrinos. Apart from its importance in realizing inflation (would be discussed in the next section), the global R-symmetry plays another important role in our analysis. Its unbroken Z_2 subgroup acts as 'matter parity', which implies a stable LSP, thereby making it a plausible candidate for dark matter. We see from Table I that baryon number violating superpotential couplings QQQ, $Q_cQ_cQ_c$ and $QQQL$ are forbidden by the $U(1)_R$ symmetry. This also holds for the higher dimensional operators, so that the proton is essentially stable [22].

Charges	S	Δ_L^a	$\bar{\Delta}_L^a$	Δ_R	$\bar{\Delta}_R$	H	L	L_c	Q	Q_c
R	2	2	0	0	0	1	0	1	0	1
Z_2	1	1	-1	1	-1	1	1	1	1	1

TABLE I: R and Z_2 charges of superfields.

III. SMOOTH HYBRID INFLATION

The superpotential term responsible for inflation is given by

$$W_{inf} = S\Big[\frac{(\Delta_R\bar{\Delta}_R)^2}{M_S^2} - M_X^2\Big]. \tag{2}$$

Note that under $U(1)_R$, S carries the same charge as W and therefore guarantees the linearity of the superpotential in S to all orders (thus excluding terms like S^2 which could ruin inflation [1].). The scalar potential derived from W_{inf} is

$$V_{inf} = \Big|\frac{(\Delta_R\bar{\Delta}_R)^2}{M_S^2} - M_X^2\Big|^2 + 4|S|^2\frac{|\Delta_R|^2|\bar{\Delta}_R|^2|}{M_S^4}\big(|\Delta_R|^2 + |\bar{\Delta}_R|^2\big) + D \text{ terms}. \tag{3}$$

Using the D-flatness condition $|\langle\Delta_R\rangle| = |\langle\bar{\Delta}_R\rangle|$, we see that the supersymmetric vacuum lies at $M = |\langle\Delta_R\rangle| = |\langle\bar{\Delta}_R\rangle| = \sqrt{M_X M_S}$ and $\langle S\rangle = 0$. Defining $\zeta/2 = |\Delta_R^0| = |\bar{\Delta}_R^0|$ and $\sigma/\sqrt{2} = |S|$, one can rewrite the scalar potential as [20, 21]

$$V_{inf} = \Big[\frac{\zeta^4}{16M_S^2} - M_X^2\Big]^2 + \frac{\sigma^2\zeta^6}{16M_S^4}. \tag{4}$$

The importance of this potential in the context of inflation is discussed in [21]. Here we can briefly summarize it. Although $\zeta = 0$ is a flat direction, it is actually a point of inflection with respect to any value of σ. It also possesses two (symmetric) valleys of local minima (containing the supersymmetric vacua) which are suitable for inflation. Unlike 'regular' supersymmetric hybrid inflation, the inclination of these valleys is already non-zero at the classical level and the end of inflation is smooth. Hence, no topological defects are produced, which is a welcome feature of smooth hybrid inflation [20, 21].

If we set $M = M_{GUT} = 2.86 \times 10^{16}$ GeV, and substitute in the expression for the quadrupole anisotropy, $(\delta T/T)_Q$, we find $M_X \simeq 1.8 \times 10^{15}$ GeV and $M_S \simeq 4.6 \times 10^{17}$ GeV [6]. Here we have employed WMAP7 [5], measurement of the amplitude of curvature perturbation ($\Delta_{\mathcal{R}}$) and set the number of e-foldings $N_Q \simeq 57$. The value of σ is 1.3×10^{17} GeV at the end of inflation (corresponding to the slow roll violating parameter, $\eta = \frac{M_P^2 V^{''}}{8\pi V} = -1$), and it is 2.7×10^{17} GeV (σ_Q) at the horizon exit. The spectral index is estimated to be $n_s \simeq 0.97$ (without supergravity corrections), close to the value of n_s from WMAP7.

Note that the supergravity corrections are important and this is studied in [6]. Once these are included (with minimal Kähler potential), n_s approaches unity (for $M \gtrsim 1.5 \times 10^{16}$ GeV) [6]. By lowering the scale M compared to the M_{GUT}, one can achieve n_s in the acceptable range. However, in this case the inflaton field-value σ_Q would be larger than the cutoff scale M_S providing a threat to the effective field theory concept.

If we employ a non-minimal Kähler potential

$$K = |S|^2 + |\Delta_R|^2 + |\bar{\Delta}_R|^2 + \frac{\kappa_S}{4}\frac{|S|^4}{M_p^2}, \tag{5}$$

then along the D-flat direction $|\Delta_R| = |\bar{\Delta}_R|$, the inflationary potential for $\sigma^2 \gg M^2$ isgiven by,

$$V = M_X^4\Big[1 - \kappa_S\frac{\sigma^2}{2M_p^2} + \big(1 - \frac{7}{2}\kappa_S + 2\kappa_S^2\big)\frac{\sigma^4}{8M_p^4} - \frac{2}{27}\frac{M^4}{\sigma^4}\Big]. \tag{6}$$

The spectral index calculated from this potential is in the desired range (0.968 ± 0.014) for different choices of κ_S. An analysis of this case is extensively studied in [8]. We have tabulated sets of values of M, M_S, σ_Q in Table II corresponding to different choices of κ_S with different predictions for the spectral index (for more examples, see Figs. 7 and 8 of [8]). With non minimal Kähler terms included, there arises the possibility of having observable tensor to scalar ratio r, a canonical measure of gravity waves produced during inflation [23].

Set	κ_S	n_s	M (GeV)	M_S(GeV)	σ_Q(GeV)
I	0	0.99	1.2×10^{16}	1.8×10^{17}	1.8×10^{17}
II	0.005	0.968	2.2×10^{16}	5.5×10^{17}	2.1×10^{17}
III	0.01	0.968	4×10^{16}	1.5×10^{18}	3×10^{17}

TABLE II: For a given value of κ_S, the predicted values of the spectral index (n_s), the gauge symmetry breaking scale (M), the cutoff scale (M_S), and the inflaton field at the time of horizon exit (σ_Q) are presented.

IV. REHEATING

Let us now discuss inflaton decay and reheating. The inflaton field(s) smoothly enter an era of damped oscillation about the supersymmetric vacuum. The oscillating system consists of two scalar fields S and $\theta = (\delta\theta + \delta\bar{\theta})/\sqrt{2}$ ($\delta\theta = \Delta_R^0 - M$ and $\delta\bar{\theta} = \bar{\Delta}_R^0 - M$) with a common mass $m_{inf} = 2\sqrt{2}M\frac{M^2}{M_S^2}$, which decay into a pair of left triplets ($\Delta_L^a, \bar{\Delta}_L^a$) and their fermionic partners ($\tilde{\Delta}_L^a, \tilde{\bar{\Delta}}_L^a$) respectively through the Lagrangian [see Eq.(2)]

$$L^s = \sqrt{2}\alpha_a\frac{M}{M_S}m_{inf}S^*\Delta_L^a\bar{\Delta}_L^a + h.c.\ ,\ L^\theta = \sqrt{2}\alpha_a\frac{M}{M_S}\theta\tilde{\Delta}_L^a\tilde{\bar{\Delta}}_L^a + h.c.. \tag{7}$$

The decay widths of both S and θ turn out to be

$$\Gamma_{inf} = \frac{3}{4\pi}\alpha_a^2\Big(\frac{M}{M_S}\Big)^2 m_{inf}\ = \frac{3}{4\pi}\Big(\frac{M_a}{M}\Big)^2 m_{inf}, \tag{8}$$

where M_a is the mass of the $SU(2)_L$ triplet given by $M_a = \alpha_a \frac{M^2}{M_S}$ (generated via the non-renormalizable superpotential coupling $\frac{\alpha_a}{M_S}\Delta_L^a\bar{\Delta}_L^a\Delta_R\bar{\Delta}_R$, after Δ_R, $\bar{\Delta}_R$ acquire vevs). For this decay to be kinematically allowed, $\alpha_a \lesssim \sqrt{2}\frac{M}{M_S}$. The splitting between M_1 and M_2 (i.e. between α_1 and α_2) will be important in estimating the lepton asymmetry. The decay of inflaton into right-handed neutrinos is kinematically forbidden since the latter have superheavy mass acquired from the renormalizable coupling $f_2 L_c L_c \Delta_R$, with f_2 of order unity.

The reheat temperature from the decay of the inflaton is $T_R \simeq \frac{1}{7}\sqrt{\Gamma M_P}$, where Γ represents the total decay width of the inflaton (here it is Γ_{inf}), where $M_P = 2.4\times 10^{18}$ GeV is the reduced Planck scale. Using the first set of values for M, M_S specified in Table II, one finds

$$T_R \simeq 0.12\times\alpha\left(\frac{M}{M_S}\right)^2\sqrt{MM_P}\ \text{GeV}, \tag{9}$$

where $\alpha = \sqrt{\alpha_1^2+\alpha_2^2}$. With the parameters involved in Table II (set II and III), we find $M/M_S \sim \mathcal{O}(10^{-2})$. Hence the reheat temperature is $T_R \sim \mathcal{O}(10^{10-11})$ GeV, where the constraint on α_a is taken into account ($\alpha \sim \mathcal{O}(10^{-3})$).

In its original form, the smooth hybrid inflationary scenario is somewhat constrained if the direct decay of inflaton into gravitinos is considered in the supergravity framework [24]. However, for heavy gravitinos, the constraint can be relaxed [24]. The exact calculation of the direct decay amplitude of the inflaton into gravitinos requires a better understanding of the supersymmetry breaking sector and the form of the Kähler potential (we are assuming here a non-minimal Kähler potential with unknown coefficients.) This is beyond the scope of this paper. Note that the high reheat temperature we obtained above does not pose much threat if the gravitino is sufficiently heavy [25] ($\gtrsim 60$ TeV), in which case it can decay well before the onset of nucleosynthesis. However, the problem may reappear from the LSPs (with $m_{LSP}\sim 100$ GeV) that are non-thermally produced from the decay of the gravitino. In ref [26], the relic abundance of these LSPs is given by

$$\Omega_{LSP}h^2 \sim 30\left(\frac{m_{LSP}}{100\ \text{GeV}}\right)\left(\frac{T_R}{10^{13}\ \text{GeV}}\right). \tag{10}$$

Hence $\Omega_{LSP}h^2 \sim 3\times 10^{-2}$ for $T_R \sim 10^{10}$ GeV, which is negligible [31]. In summary, we conclude that at the end of inflation, the inflaton system has decayed away into $SU(2)_L$ triplets. We will show in the next section that the subsequent decay of these $SU(2)_L$ triplets creates a lepton asymmetry, which is partially converted into the observed baryon asymmetry via the electroweak sphaleron effects [28].

V. TYPE II NON-THERMAL LEPTOGENESIS AND NEUTRINO MASSES

In general both the right-handed neutrinos as well as the left-handed triplets can yield a lepton asymmetry in left-right models [18]. However, in our case with superheavy ($M\sim 10^{16}$ GeV) right handed neutrinos, the Leptogenesis would come mainly from the $SU(2)_L$ triplets. Note that we have considered two pairs of $SU(2)_L$ superfields, so that the CP asymmetry would be nonzero.

The experimental value of the baryon to photon ratio is given by [5]

$$\frac{n_B}{n_\gamma} \simeq (6.5\pm 0.4)\times 10^{-10}. \tag{11}$$

In this respect, the required lepton asymmetry is estimated to be

$$\left|\frac{n_L}{s}\right| \simeq (2.67-3.02)\times 10^{-10}. \tag{12}$$

To estimate the lepton asymmetry we follow the analysis of ref [16]. The Higgs triplet Δ_L^a decays into LL and HH (see Fig. 1(a)), while $\bar{\Delta}_L^a$ decays into $\tilde{L}\tilde{L}$ and $\tilde{H}\tilde{H}$. The amount of CP violation in these decays is controlled by the interference of the tree level process with one-loop diagram (see Fig. 1(b)) as described in [16].

The effective mass-squared matrix of the scalar triplets Δ_L^a and $\bar{\Delta}_L^a$ is [16], $\Delta_L^{a\dagger}(M^2)_{ab}\Delta_L^b + \bar{\Delta}_L^{a\dagger}(M'^2)_{ab}\bar{\Delta}_L^b$, where

$$M^2 = \begin{pmatrix} M_1^2 - i\Gamma_{11}M_1 & -i\Gamma_{12}M_2 \\ -i\Gamma_{21}M_1 & M_2^2 - i\Gamma_{22}M_2 \end{pmatrix}, \tag{13}$$

and M'^2 has a similar pattern with Γ_{ab} replaced by Γ'_{ab}. The contributions to Γ_{ab} (Γ'_{ab}) come from the absorptive part of the one loop self-energy diagrams for $\Delta_L^a \to \Delta_L^b$ ($\bar{\Delta}_L^a \to \bar{\Delta}_L^b$),

$$\Gamma_{ab}M_b = \frac{1}{8\pi}[\Sigma_{ij}(f_{1ij}^{a*}f_{1ij}^b)p_{\Delta_L}^2 + M_aM_bg^ag^{b*}]\ ,$$

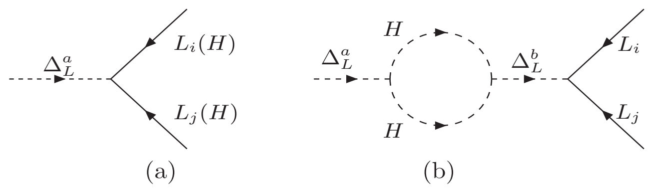

FIG. 1: (a) Tree level decay(s) of Δ_L into leptons (Higgs). (b) One loop self energy diagram for the generation of CP asymmetry.

$$\Gamma'_{ab} M_b = \frac{1}{8\pi}[\Sigma_{ij}(f^a_{1ij} f^{b*}_{1ij}) M_a M_b + p^2_{\Delta_L} g^{a*} g^b], \tag{14}$$

where i, j are generation indices, $g^a = \gamma^a(\frac{M}{M_S})$ and p^2_Δ is the momentum squared of the incoming or outgoing particle. The physical states $\chi^{1,2}_+$, $\xi^{1,2}_+$ (with masses $\sim M_{1,2}$) can be obtained[32] by diagonalizing M^2, M'^2. Here we neglect terms of order $[\frac{\Gamma_{ij} M_j}{M_1^2 - M_2^2}]^2$.

The CP asymmetries are then defined by [16]

$$\begin{aligned} \epsilon^a &= \triangle L \frac{\Gamma(\chi^a_- \to ll) - \Gamma(\chi^a_+ \to l^c l^c)}{\Gamma_{\chi^a_-} + \Gamma_{\chi^a_+}}, \\ &= \frac{M_1 M_2}{2\pi(M_1^2 - M_2^2)} \frac{\Sigma_{ij} \mathrm{Im} f^1_{1ij} f^{2*}_{1ij} g^1 g^{2*}}{\Sigma_{ij} |f^a_{1ij}|^2 + |g^a|^2}, \end{aligned} \tag{15}$$

and

$$\begin{aligned} \epsilon'^a &= \triangle L \frac{\Gamma(\xi^a_+ \to ll) - \Gamma(\xi^a_- \to l^c l^c)}{\Gamma_{\xi^a_+} + \Gamma_{\xi^a_-}}, \\ &= \frac{M_1 M_2}{2\pi(M_1^2 - M_2^2)} \frac{\Sigma_{ij} \mathrm{Im} f^1_{1ij} f^{2*}_{1ij} g^1 g^{2*}}{\Sigma_{ij} |f^a_{1ij}|^2 + |g^a|^2}, \end{aligned} \tag{16}$$

where the lepton number violation ΔL changes by 2 units. We note that $\epsilon^a = \epsilon'^a$.

The lepton asymmetry is given by

$$\begin{aligned} \frac{n_L}{s} &\simeq \frac{3}{2} \frac{T_R}{m_{inf}} \Sigma_a 3[\epsilon^a + \epsilon'^a], \\ &= \Sigma_a \frac{3}{2} \frac{T_R}{m_{inf}} \frac{3 M_1 M_2}{\pi(M_1^2 - M_2^2)} \frac{\Sigma_{ij} \mathrm{Im} f^1_{1ij} f^{2*}_{1ij} g^1 g^{2*}}{\Sigma_{ij} |f^a_{1ij}|^2 + |g^a|^2}, \end{aligned} \tag{17}$$

where the ratio of the number density of the $SU(2)_L$ triplets (n_Δ) to the entropy density s is expressed as $\frac{3}{2}\frac{T_R}{m_{inf}}$. Once this asymmetry is created, one should ensure that it is not erased by the lepton-number non-conserving interactions (for example $HH \leftarrow \Delta_L \to LL$, $\tilde{H}\tilde{H} \leftarrow \bar{\Delta}_L \to \tilde{L}\tilde{L}$). As long as the $SU(2)_L$ triplet masses (M_a) are sufficiently larger than T_R (here $\frac{T_R}{M_a} \simeq 0.12 \frac{\sqrt{M M_P}}{M_S}$ with the specific choice of M, M_S as given in Table II (set II and III), there will be no significant wash-out factor, unlike thermal leptogenesis.

To estimate n_L/s, we need to fix some parameters appearing in Eq.(17) which are also involved in the light neutrino mass matrix. The neutrino mass matrix is represented by the type II see-saw relation

$$m_\nu = 2 f^a_{1ij} v^a_{\Delta_L} - m_D^T M_R^{-1} m_D \equiv m_{\nu_{II}} - m_{\nu_I}, \tag{18}$$

where $v^a_{\Delta_L}$ are the $SU(2)_L$ triplet Higgs's vevs. With the masses of all right handed neutrinos comparable to M, m_{ν_I} are too small to account for the solar and atmospheric neutrino data. Hence $m_{\nu_{II}}$ provides the main contribution to

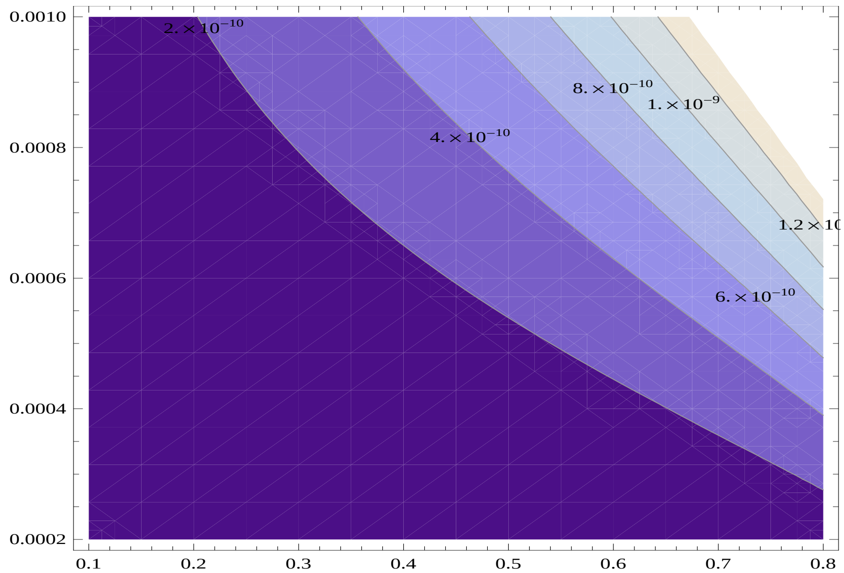

FIG. 2: Contour plot for n_L/s as a function of the parameters: $p = M_2/M_1$ and $g \lesssim M/M_S$.

the neutrino mass matrix, namely

$$(m_\nu)_{ij} \simeq 2 f^a_{1ij} \frac{g^a}{M_a} v^2, \tag{19}$$

where $v \simeq 174$ GeV. In order to estimate both the lepton asymmetry (Eq.(17)) and neutrino masses (through Eq.(19)), we first simplify by assuming $|g^1| \simeq |g^2| = g$, $|f^1_1| \simeq |f^2_1| = f_1$ (thus $|\Sigma_{ij} f^1_{1ij} f^{2*}_{1ij}| \simeq \Sigma_{ij}|f_{1ij}|^2$). Then Eqs.(17) and (19) can be expressed as

$$\begin{aligned}\frac{n_L}{s} &\simeq \frac{9}{\pi}\frac{T_R}{m_{inf}} \times \frac{M_1 M_2}{M_1^2 - M_2^2} \times \frac{\Sigma_{ij}|f_{1ij}|^2 g^2}{\Sigma_{ij}|f_{1ij}|^2 + g^2}, \\ &\simeq \frac{0.374}{\pi}\sqrt{\frac{M_P}{M}}\alpha \times \frac{M_1 M_2}{M_1^2 - M_2^2} \times \frac{\Sigma_{ij}|f_{1ij}|^2 g^2}{\Sigma_{ij}|f_{1ij}|^2 + g^2}, \qquad (20)\\ (m_\nu)_{ij} &\simeq 2 f_{1ij} g v^2 \Big(\frac{1}{M_1} + \frac{1}{M_2}\Big), \qquad (21)\end{aligned}$$

where we have substituted for T_R and M_a and assumed the CP violating phase to be maximal.

The neutrino mass matrix m_ν can be diagonalized by

$$m_\nu = U^*_\nu m^{diag}_\nu U^\dagger_\nu, \tag{22}$$

where $m^{diag}_\nu = \mathrm{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$. In the basis where the charged lepton matrix is diagonal, U_ν coincides with the lepton mixing matrix. Using Eqs.(21), we get

$$\frac{n_L}{s} \simeq \frac{0.374}{\pi}\frac{p\sqrt{1+p^2}}{1-p^2}\sqrt{\frac{M_P}{M}}\frac{M_1}{M}\frac{M_S}{M} \times \frac{\Sigma_{ij}|m_{\nu_{ij}}|^2 F g^2}{\Sigma_{ij}|m_{\nu_{ij}}|^2 F + g^4}, \tag{23}$$

where $F = \frac{M_1^2 M_2^2}{4v^4(M_1+M_2)^2} = \frac{p^2}{(1+p)^2} \times \frac{M_1^2}{4v^4}$. Here p determines the degree of degeneracy between M_1 and M_2, defined by $M_2 = pM_1$. Since the parameter g is defined as $g^a = \gamma^a \frac{M}{M_S}$, its maximum value is of order $\frac{M}{M_S}$. Finally, using the current experimental limits for neutrino masses [29], one finds that $\Sigma_{ij}|m_{\nu_{ij}}|^2$ is given by $\Sigma_{ij}|m_{\nu_{ij}}|^2 \simeq 0.0025$ eV^2, where we have used the best fitted values of the neutrino mixing angles $\theta_{12}, \theta_{23}, \theta_{13}$ and mass squared differences [29]. We have taken the lightest neutrino mass eigenvalue to be zero. In Fig. 2 we present the lepton asymmetry as a function of p and g with $\alpha_1 = 10^{-3}$. We see that n_L/s can be of order the desired value $(2-3)\times 10^{-10}$ for $0.2 \lesssim p \lesssim 0.8$ and $g \gtrsim 2.5 \times 10^{-4}$, which means $\gamma^a \simeq \mathcal{O}(0.01)$. It is worth mentioning that with these values one finds that M_1 and M_2 are given by $M_1 \simeq 10^{12}$ GeV and $M_2 \simeq (2-8)\times 10^{11}$ GeV. Therefore, $M_{1,2}/T_R > 10$, which indicates that no washout should happen.

VI. CONCLUSIONS

We have considered type II non-thermal leptogenesis in the context of smooth hybrid inflation. The scheme is consistent with the observed solar and atmospheric neutrino oscillations. Although our discussion is based on the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, it is clear that it could be extended to other models which contain suitable $SU(2)_L$ triplet scalars with tiny vevs responsible for the observed neutrino masses. The stability of the proton will depend on the underlying gauge symmetry.

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[31] For studies on the constraints on the reheating temperature in general, see [27].
[32] The physical states $\chi_-^{1,2}$ ($\xi_-^{1,2}$) are also similarly obtained by diagonalizing the matrix in Eq.(13) with Γ_{12} replaced by Γ_{12}^* and vice versa.