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No-scale gauge non-singlet inflation inducing TeV scale inverse seesaw mechanism

Ahmad Moursy

*Department of Basic Sciences, Faculty of Computers and Artificial Intelligence,
Cairo University, Giza 12613, Egypt*

E-mail: a.moursy@fci-cu.edu.eg

ABSTRACT: We develop a model of sneutrino inflation that is charged under $U(1)_{B-L}$ gauge symmetry, in no-scale supergravity framework. The model provides an interesting modification of tribid inflation. We impose $U(1)_R$ symmetry on the renormalizable level while allow Planck suppressed non-renormalizable operators that break R-symmetry. This plays a crucial role in realizing a Starobinsky-like inflation scenario from one hand. On the other hand it plays an essential role, as well as SUSY breaking effects, in deriving the tiny neutrino masses via TeV inverse seesaw mechanism. Thus, we provide an interpretation for the extremely small value of the μ_S mass parameter required for inverse seesaw mechanism. We discuss a reheating scenario and possible constraints on the model parameter space in connection to neutrino masses.

KEYWORDS: Beyond Standard Model, Cosmology of Theories beyond the SM, Supergravity Models, Supersymmetry Breaking

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Contents

1	Introduction	1
2	The model	2
3	Inflation scenario and observables	5
3.1	Inflation trajectory and effective potential	5
3.2	Inflation observables	8
3.3	Loop corrections	10
4	SUSY breaking, reheating and neutrino masses	11
4.1	Constraints from reheating scenarios	12
4.2	Neutrino masses	14
5	Conclusions	16

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1 Introduction

Remarkable studies over the past few decades were devoted to explore the connection and interplay between different hierarchical scales in nature. Finding such connections helps in understanding the fundamental origin of these scales. This issue has attracted careful studies on the huge hierarchy between the scales at the low energy observable sector, such as electroweak and neutrino mass scales from one hand, and on the other hand the early universe scales such as Planck, GUT, SUSY breaking and inflation scales.

Inflation paradigm is supported by CMB observations. Planck collaboration [1] has confirmed the values of the spectral index of the scalar fluctuations $n_s = 0.955 - 0.974$, up to 2 sigma exclusion limits, and the tensor to scalar ratio $r < 0.08$. One of the appealing models of inflation is the Starobinsky model [7–9] whose predictions are consistent with Planck observations. Starobinsky-like potentials can be implemented in supergravity framework such as no-scale supergravity models [10–16]. The latter models have a striking feature of solving naturally the η -problem without any need to preserve R-symmetry that is required in supersymmetric hybrid inflation models [19, 21].

It is tempting to study the link between inflation scale and neutrino mass scale. The detected neutrino oscillations [2–6] implied that neutrinos have tiny masses. The latter hinted at new physics beyond the standard model that accounts for seesaw mechanisms [27–30]. Interestingly right-handed neutrino supermultiplets can be employed to play a double role in realizing seesaw type I mechanism using its fermionic components, while the scalar component (right-handed sneutrino) can be assigned as an inflaton, with mass scale $\sim 10^{13}$ GeV. Chaotic inflation with a singlet right-handed sneutrino was proposed in [20].

In [21] sneutrino hybrid inflation was introduced in supergravity framework with exact R-symmetry while Heisenberg symmetry was imposed in [22, 23], in order to solve the η -problem. Non-singlet sneutrino inflation was explored in [18, 24], where the η -problem was avoided by either Heisenberg symmetry or shift symmetry.

In this work we consider a non-singlet inflaton under $U(1)_{B-L}$ gauge symmetry that is assigned to right-handed sneutrino, in no-scale supergravity. One remarkable feature of the present model is that the $U(1)_{B-L}$ gauge symmetry is broken during inflation by the inflaton hence topological defects such as cosmic strings are absent. After the inflation ends, the gauge symmetry is broken in the true minimum at GUT scale which is a derived scale from Planck scale M_P and inflationary mass scale μ . Moreover, the inflation fields play a crucial rule in generating tiny neutrino masses via a TeV inverse seesaw mechanism [31–33] that allows non-small values of the neutrino Yukawa couplings Y_ν , without unnatural fine tuning. We introduce $U(1)_R$ symmetry (R-symmetry) on the renormalizable level while we allow non-renormalizable operators suppressed by Planck mass to break R-symmetry [15, 17].

The paper is organized as follows. In section 2 we construct the model discussing the role of each supermultiplet. In section 3 we explore the inflation scenario. Also, we analyze the inflation effective potential, trajectory and observables. We discuss consequences on post inflation era in section 4. We point out an interesting feature of SUSY breaking in models of hybrid inflation and exist in our model. We investigate constraints that arise from reheating and then generate neutrino masses from inverse seesaw mechanism. Finally, we draw our conclusions in section 5.

2 The model

We consider a no-scale supergravity model in which both the superpotential and the Kähler potential are invariant under $U(1)_{B-L}$ gauge group. Moreover, on the renormalizable level, the Lagrangian respects the $U(1)_R$ symmetry, with $R[W] = 2$ while non-renormalizable operators suppressed by Planck mass are allowed to break R-symmetry [15, 17]. A Z_3 discrete symmetry is imposed in order to guarantee a successful inflation and TeV inverse seesaw mechanism. The field content of the model is listed as follows.

- The supermultiplets $S_1 = (\tilde{S}_1, S_1)$, $S_2 = (\tilde{S}_2, S_2)$ are gauge non-singlets under $U(1)_{B-L}$. Their scalar components are corresponding to the inflaton while the fermionic components mix with the right-handed neutrino N in the neutrino mass matrix.
- The waterfall supermultiplets $\phi_1 = (\phi_1, \tilde{\phi}_1)$, $\phi_2 = (\phi_2, \tilde{\phi}_2)$ whose scalar components are frozen at the origin during the inflation and after the inflation ends they break the $B - L$ gauge symmetry at the true minimum.
- The singlet $S = (S, \tilde{S})$ plays a five-fold role: first, it is a driving field where its F-term generates the potential of the waterfall fields, (the scalar components of ϕ_1, ϕ_2), hence they acquire vevs to break the $B - L$ symmetry at GUT scale $M \sim \sqrt{\mu M_P}$. Second,

	S_1	S_2	ϕ_1	ϕ_2	S	N
$U(1)_{B-L}$	1	-1	2	-2	0	-1
R	1	1	0	0	2	-1
Z_3	ω^2	ω	ω	ω^2	1	ω

Table 1. $U(1)_{B-L}$ and R charge assignments for different superfields $S_1, S_2, \phi_1, \phi_2, S, N$ in the inflation sector.

the non-renormalizable coupling to S_1, S_2 will play important role in flattening the inflation potential causing wider region of the parameter space of a successful inflation as we will demonstrate in the following section. Third, its decay contributes to reheating the universe. On the other hand SUSY breaking effects generate a tadpole term in the scalar S shifting its minimum from the origin, hence the fourth role is realized in generating the mixing term $\mu_S S_1 S_2$ after the end of inflation and can contribute to the neutrino mass matrix. Finally non-zero vev of S can be employed to solve the μ -problem of the TeV scale MSSM by generating the TeV scale mixing term $\mu H_u H_d$ that is needed to trigger the electroweak symmetry breaking [38].

- The right-handed neutrino supermultiplet $N = (\tilde{N}, N)$ which contributes to the neutrino mass matrix and has no role in the inflation scenario.

The superfields constituting the model, which are SM singlets, as well as the respective $B - L$, the R-symmetry and Z_3 charges are given in table 1. In this respect the superpotential and the Kähler potential for the inflation sector are given by

$$W_{\text{inf}} = \kappa_1 S (\phi_1 \phi_2 - \mu M_P) + \kappa_2 S_1 S_2 \left(\frac{\phi_1 \phi_2}{M_P} - \mu \right) + \frac{\lambda_1}{M_P} (S_1 S_2)^2 + \frac{\lambda_2}{M_P} S^2 S_1 S_2, \quad (2.1)$$

$$K = -3M_P^2 \log \left[\frac{T + \bar{T}}{M_P} - \frac{|S|^2}{3M_P^2} - \frac{|S_1|^2}{3M_P^2} - \frac{|S_2|^2}{3M_P^2} - \frac{|\phi_1|^2}{3M_P^2} - \frac{|\phi_2|^2}{3M_P^2} \right], \quad (2.2)$$

The superpotential (2.1) is a development of the tribrid inflation model [21–26]. As will be detailed later, the difference is at two main points: the first is that the mass term of right-handed neutrinos $S_1 S_2$ vanishes at the end of inflation. The second is that the vev of S is not zero during the inflation. The parameter μ determines the scale of the inflation and κ_i, λ_i are dimensionless couplings where λ_i are responsible for the R-symmetry breaking in the superpotential and $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass.¹ In this regard, the total scalar potential consists of the F-term and D-term parts $V = V_F + V_D$. The F-term scalar potential is given by

$$V_F = e^{K/M_P^2} \left[D_I K^{I\bar{J}} D_{\bar{J}} \bar{W} - 3 \frac{|W|^2}{M_P^2} \right], \quad (2.3)$$

¹Worthy to mention that the symmetry allows the Planck suppressed non-renormalizable operators such as S^4 , $S^2 \phi_1 \phi_2$ and $(\phi_1 \phi_2)^2$. However they are undesirable hence we can set the corresponding couplings to zero or we can argue that they are forbidden in the UV sector by R -symmetry. We assume that only S, S_i, N as well as the MSSM superfields carry R-charges. We are not going to give more details on the UV sector as it is beyond the scope of the paper.

where I, J run over $(T, S, S_1, S_2, \phi_1, \phi_2)$, $K^{I\bar{J}}$ is the inverse of the Kähler metric $K_{I\bar{J}} = \frac{\partial K}{\partial Z^I \partial \bar{Z}^{\bar{J}}}$, and D_I is the Kähler derivative defined by $D_I = \frac{\partial}{\partial Z^I} + \frac{1}{M_P^2} \frac{\partial K}{\partial Z^I}$. We use the lower case letters i, j to run over the inflation sector fields $S, S_1, S_2, \phi_1, \phi_2$. The D-term potential is given by

$$V_D = \frac{g^2}{2} \text{Re} f_{AB}^{-1} D^A D^B, \quad (2.4)$$

where f_{AB} is the gauge Kinetic function and the indices A, B are corresponding to a representation of the gauge group under which Z^i are charged. The D-term D^A is given by

$$D^A = \frac{\partial K}{\partial Z^i} (T^A)^i_j Z^j, \quad (2.5)$$

with T^A are generators of the gauge group in the appropriate representation. We will work in the D-flat direction, hence the total potential will be given by F-term scalar potential as

$$V = \frac{1}{\Omega^2} \sum_{i=1}^5 \left| \frac{\partial W}{\partial Z_i} \right|^2, \quad \Omega = \frac{T + \bar{T}}{M_P} - \frac{|S|^2}{3M_P^2} - \frac{|\tilde{S}_1|^2}{3M_P^2} - \frac{|\tilde{S}_2|^2}{3M_P^2} - \frac{|\phi_1|^2}{3M_P^2} - \frac{|\phi_2|^2}{3M_P^2}. \quad (2.6)$$

After calculating the scalar potential, we consider the modulus T to be stabilized at high scale by any of the mechanisms of [13, 34, 35], where $\langle \text{Re}(T) \rangle = \tau_0$, $\langle \text{Im}(T) \rangle = 0$. It is clear that, the total scalar potential is positive semidefinite. Therefore it has a global SUSY minimum which is Minkowskian, and is located at

$$\langle S \rangle = \langle \tilde{S}_1 \rangle = \langle \tilde{S}_2 \rangle = 0 \text{ \& } \langle \phi_1 \phi_2 \rangle = \mu M_P \quad (2.7)$$

We will work in the D-flat direction associated with the $B - L$ symmetry,

$$|\tilde{S}_1|^2 - |\tilde{S}_2|^2 + 2|\phi_1|^2 - 2|\phi_2|^2 = 0. \quad (2.8)$$

Without loss of generality we choose for simplicity the direction $|\tilde{S}_1| = |\tilde{S}_2| = |\Sigma|$ and $|\phi_1| = |\phi_2| = |\phi|$, that cancels two real degrees of freedom. Therefore, the scalar potential simplifies to the following form

$$V = \frac{1}{\Omega^2} \left[\left| \kappa_1 (\phi^2 - \mu M_P) + \frac{2\lambda_2}{M_P} S \Sigma^2 \right|^2 + 2|\phi|^2 \left| \kappa_1 S + \frac{\kappa_2}{M_P} \Sigma^2 \right|^2 + 2|\Sigma|^2 \left| \frac{\kappa_2}{M_P} (\phi^2 - \mu M_P) + \frac{\lambda_1}{M_P} \Sigma^2 + \frac{\lambda_2}{M_P} S^2 \right|^2 \right]. \quad (2.9)$$

The previous direction is not unique as one can choose another D-flat direction in which $|\phi_1|^2 - |\phi_2|^2 = \gamma \neq 0$, which implies that $|\tilde{S}_1|^2 = |\tilde{S}_2|^2 - 2\gamma$, hence only one degree of freedom cancels. However, as we will show in the next section, all the four real degrees of freedom $\alpha_{1,2}, \beta_{1,2}$ are fixed to zero during inflation, where $\phi_j = \frac{\alpha_j + i\beta_j}{\sqrt{2}}$. Hence, during the inflation and at the global minimum $\gamma = 0$, and we return to the simple case of $|\tilde{S}_1| = |\tilde{S}_2| = |\Sigma|$. In this regard we express the complex fields S, Σ, ϕ in terms of their real components as follows

$$S = \frac{s + i\sigma}{\sqrt{2}}, \quad \phi = \frac{\alpha + i\beta}{\sqrt{2}}, \quad \Sigma = \frac{\chi + i\psi}{\sqrt{2}}. \quad (2.10)$$

Figure 1. The scalar potential of x, α near the SUSY vacuum (left panel) and for large values of x (right panel), with $\beta = y = \sigma = 0$ and $\lambda_1 = 0.3 \times 10^{-5}, \mu = \times 10^{-5}, \kappa_1 = 10^{-4}, \kappa_2 \sim 1$. All values are given in the units where $M_P = 1$.

Accordingly, the global minimum of the potential is located at

$$s = \sigma = \chi = \psi = \beta = 0 \ \& \ \alpha = \sqrt{2\mu M_P}. \quad (2.11)$$

It turns out that χ and ψ are massless at the true minimum, since the second term of (2.1) vanishes when the higgs fields acquire non-zero vev $\sim \sqrt{\mu M_P}$. However, as we will demonstrate later, SUSY breaking effects will generate mass terms for $\tilde{S}_i$. On the other hand the squared masses of the other real components are given by

$$m_s^2 = m_\sigma^2 = \frac{3\kappa_1^2 \mu M_P^2}{3\tau_0 - \mu} \quad \text{and} \quad m_\alpha^2 = m_\beta^2 = \frac{\kappa_1^2 \mu M_P^2}{\tau_0}. \quad (2.12)$$

3 Inflation scenario and observables

3.1 Inflation trajectory and effective potential

Here we study the trajectory of inflation and the resulting effective inflation potential. The scalar potential is minimized at $\phi_1 = \phi_2 = 0$ for large values of S_1, S_2 . Along the inflation trajectory $\phi_1 = \phi_2 = 0$ and the effective inflation potential receives a dominant contribution from the F-terms F_{S_i} as well as a contribution from F_S , that will be constrained from the inflation observables, while the contribution of F_{ϕ_i} vanishes. On the other hand, the inflation energy shifts the vev of s from zero during inflation to a small value, due to the coupling $\frac{\lambda_2}{M_P} S^2 S_1 S_2$. We depict the shape of the scalar potential of (x, α) during the inflation and near the global SUSY minimum after the inflation ends in figure 1.

We first investigate the form of the scalar potential in the limit $\lambda_2 = 0$ (the potential is minimized at $S = 0$), then we expand the scalar potential for small s when $\lambda_2 \neq 0$. In this respect, the inflation potential, when $\lambda_2 = 0$, has the form

$$V(S=0) = \frac{9M_P^2 \left[8\lambda_1^2 |\Sigma|^6 + \kappa_1^2 \mu^2 M_P^4 + 2\kappa_2^2 \mu^2 M_P^2 |\Sigma|^2 - 4\kappa_2 \lambda_1 \mu M_P (\bar{\Sigma}^4 + \Sigma^4) \right]}{4(|\Sigma|^2 - 3M_P \tau_0)^2}. \quad (3.1)$$

In order to have canonical kinetic terms we use the following field redefinitions

$$\Sigma = \sqrt{3M_P\tau_0} \tanh\left(\frac{x+iy}{\sqrt{6}M_P}\right), \quad (3.2)$$

where x, y are real scalar fields. As will be demonstrated latter, the field y acquires a mass of order Hubble scale during the inflation hence it is stabilized at zero [16]. Therefore, the scalar potential takes the form

$$V = A \cosh^4\left(\frac{x}{\sqrt{6}}\right) \left[B^2 \tanh^6\left(\frac{x}{\sqrt{6}}\right) - 2B \tanh^4\left(\frac{x}{\sqrt{6}}\right) + \tanh^2\left(\frac{x}{\sqrt{6}}\right) + f \right], \quad (3.3)$$

where,

$$A = \frac{3\mu^2\kappa_2^2}{2\tau_0 M_P}, \quad B = \frac{6\lambda_1\tau_0}{\kappa_2\mu}, \quad f = \frac{\kappa_1^2 M_P}{6\kappa_2^2\tau_0}, \quad (3.4)$$

are dimensionless parameters that will be determined by slow-roll conditions and inflation observables. We will use the units in which $M_P = 1$ throughout this section. When $f = 0$ (or equivalently $\kappa_1 = 0$), we obtain the potential studied in [15]. Therefore, in the limit where $B = 1$ and $f = 0$, the potential (3.3) is asymptotically flat

$$V_0 = A \tanh^2\left(\frac{x}{\sqrt{6}}\right). \quad (3.5)$$

Now we return to consider the effect of the coupling $\frac{\lambda_2}{M_P} S^2 S_1 S_2$. It gives rise to a linear term in s in the F_S contribution to the potential (2.9), hence s acquires small vev during the inflation and it is given by

$$s = \frac{\sqrt{2A}f}{3\lambda_2}, \quad (3.6)$$

where we have expanded for large x during the inflation and taken the leading order.²

Now, we consider for completeness, that all real components of $\phi_{1,2}$ that can account for possible D-flat directions, α_i, β_i . We show that they are fixed at zero during inflation as well as the real degrees of freedom y, σ . As a matter of fact, all real degrees of freedom, apart from the inflaton x , acquire masses larger than the Hubble scale H during inflation [16]. We present the ratio of field dependent squared masses to the squared Hubble scale during inflation in figure 2.

Returning to the scalar α that accounts for the waterfall phase, the scalar potential (2.9) is minimized at $\alpha = 0$ during the inflation, for x is rolling at large values, since the field dependent squared mass m_α^2 is positive and is of order squared Hubble scale during inflation. When x reaches a critical value x_c after the inflation ends, m_α^2 flips its sign to negative values which triggers the waterfall phase [16]. Expanding for small values of the inflaton x , we have

$$m_\alpha^2 \simeq \frac{\kappa_1^2 \mu (\mu - 3\tau_0)}{6\tau_0^2}, \quad (3.7)$$

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²This will not affect the canonical kinetic term of the inflaton as s is much small during inflation.

Figure 2. The ratio of the squared mass of the real scalar fields α_i, β_i, y and σ , to the squared Hubble scale, with $A = 3 \times 10^{-10}$, $f = 3.33 \times 10^{-11}$, $\lambda_2 = 0.005$, $B = 1$. All values are given in the units where $M_P = 1$.

Figure 3. The scalar potential $\frac{V_{\text{inf}}}{A}$ of the inflaton x in the two cases when s is zero (solid curves) and non-zero (dashed curves), with $A = 3 \times 10^{-10}$, $\lambda_2 = 0.5$, $B = 1$ and we took two values for $f = 3.33 \times 10^{-9}$ and 3.33×10^{-5} . All values are given in the units where $M_P = 1$.

which is negative as long as $\mu < 3\tau_0$. The critical value of the inflaton x_c will be given as

$$x_c = \frac{2\kappa_2\mu}{\tau_0(\kappa_2 + 4\lambda_1)}. \quad (3.8)$$

We expand the scalar potential (2.9) to second order in s with the other fields being fixed at the origin, and consider the limit where $B = 1$, then we obtain the inflation effective potential as follows

$$V_{\text{inf}} = V_0 + A f \cosh^4\left(\frac{x}{\sqrt{6}}\right) - 6\lambda_2 \sqrt{\frac{f}{2A}} s V_0 \cosh^4\left(\frac{x}{\sqrt{6}}\right) + \frac{s^2 \cosh^2\left(\frac{x}{\sqrt{6}}\right)}{6A\tau_0} \\ \times \left[A f \cosh^4\left(\frac{x}{\sqrt{6}}\right) + V_0 \left(-3A\lambda_2 \sqrt{\frac{6\tau_0}{A}} + A + 27\lambda_2^2 \tau_0 \sinh^2\left(\frac{x}{\sqrt{6}}\right) \right) \right]. \quad (3.9)$$

Figure 4. The inflation potential $\frac{V_{\text{inf}}}{A}$ with $A = 3 \times 10^{-10}$, $\lambda_2 = 0.5$, $s \neq 0$ and $s = 0$ for the solid and dashed curves respectively, and we took two values for $f = 3.33 \times 10^{-9}$ and 3.33×10^{-5} . The left panel corresponds to $B = 1 + 10^{-5}$ and the right panel corresponds to $B = 1 - 10^{-5}$. All values are given in the units where $M_P = 1$.

If $\lambda_2 = 0$, then s will be stabilized at zero during the inflation, hence the third and fourth terms will be absent in the inflation potential (3.9), and we return to the potential (3.3). It is clear that the second term proportional to f spoils the asymptotic flatness of the potential at large values of the inflaton x , hence f should be small enough in order to guarantee the flatness of the potential during the inflation. For small values of s , the third term dominates over the fourth term during the inflation. As a consequence, the third term in (3.9) has a desirable effect in cancelling the dangerous effect of the second term hence flattening the potential during the inflation. In figure 3 we show the inflation potential for $B = 1$ when s is stabilized at zero and non-zero values during inflation where it was indicated that the potential is flattened for $s \neq 0$ during the inflation as depicted in the solid curves. When $B \neq 1$, we define $B(\xi \neq 0) = 1 + \xi$, hence extra terms proportional to ξ are generated and this may spoil the flatness during the inflation. Accordingly the constraint $|\xi| \lesssim \mathcal{O}(10^{-4})$ should be imposed [15]. Figure 4 illustrates the inflation potential for $B \neq 1$ ($\xi \neq 0$) for different values of f . The solid curves corresponding to non-zero values of s , which coincide for different values of f .

3.2 Inflation observables

Now we move to discuss the model predictions of inflation observables and explore the possible constraints on the different parameters. The inflation observables, tensor-to-scalar ratio r , spectral tilt n_s and the scalar amplitude A_s will be investigated. In terms of the slow-roll parameters ϵ and η , the observables are defined as follows

$$\begin{aligned} r &= 16\epsilon, \\ n_s &= 1 - 6\epsilon + 2\eta, \\ A_s &= \frac{V}{24\pi^2\epsilon}, \end{aligned}$$

where the above observables are computed at the crossing horizon value of the inflaton field x_* . The number of e-folds is given by

$$N = \int_{x_e}^{x_*} \frac{1}{\sqrt{2\epsilon}} dx, \quad (3.10)$$

Figure 5. A logarithmic plot for the model predictions in the n_s – r plane. We have fixed $A \sim 10^{-10}$. The solid (dash-dotted) blue curves correspond to $N = 60$ ($N = 50$) respectively where we scanned over $B = 1 + \xi$ with $\xi \in [-10^{-4}, 10^{-4}]$, with $f = \lambda_2 = 0$. The other segments correspond to fixed $B = 1$ ($\xi = 0$) with different values of f , λ_2 and we have scanned over $N = 50 - 60$ (small-big) dots, for each segment. The dark and light red regions correspond to the 1 and 2 sigma exclusion limits released by the Planck collaboration (2018) (TT+ TE+ EE + LowE + Lensing + BK14) [1].

where x_e is the value of the inflaton at the end of inflation. The observed value of the scalar amplitude $A_s \simeq 1.95896 \pm 0.10576 \times 10^{-9}$ at 68% CL [1], which fixes $A \sim 10^{-10}$.

In figure 5, we displayed the predictions of our model for n_s , r in a logarithmic plot versus the Planck limits. First, we fixed $f = \lambda_2 = 0$ and scanned over $B = 1 + \xi$ with $-10^{-4} \leq \xi \leq 10^{-4}$ on the solid (dash-dotted) blue curves for $N = 60$ ($N = 50$), respectively. On the other segments, we scanned over $N = 50 - 60$, and fix $B = 1$ ($\xi = 0$). The orange segment illustrates the situation when the potential is asymptotically flat (3.5) where $f = \lambda_2 = 0$. However, f (or κ_1) should not be zero in order to keep the potential terms responsible for stabilizing the waterfall fields at the true minimum. Therefore fixing $\lambda_2 = 0$ (which means that $s = 0$ on the inflation trajectory), the inflation observables constrain $f \lesssim 2 \times 10^{-6}$ as indicated in the green dashed segment. Turning on λ_2 to non-zero value will flatten the potential causing successful inflation. The latter case is illustrated in the solid green segment where we took the same value of f on the dashed one and let $\lambda_2 \sim 10^{-6}$. Any greater value of λ_2 makes the potential more flat at larger inflaton value. Therefore we investigate greater values of $f \lesssim \mathcal{O}(1)$ that spoil the flatness of the potential against greater values of λ_2 . The Purple segment corresponds to $\lambda_2 \sim 0.05$ and $f \sim 1$. For larger value of λ_2 purple segment moves towards the orange one. For $\lambda_2 \gtrsim 0.5$, the purple and orange segments coincide. For $B \neq 1$ and $\lambda_2 \gtrsim 0.5$, different potentials with different values of f coincide as depicted in figure (4). In order to have successful inflation, the constraint $|\xi| \lesssim \mathcal{O}(10^{-4})$ should be satisfied. In table 2, we have provided some specific values of the superpotential parameters as well as the modulus vev τ_0 corresponding to the values of A, f, λ_2 and $B = 1$, that were discussed above in the inflation observables. We have focused on showing different regimes of small and large values of κ_1 and κ_2 , since they

λ_2	f	μ	κ_1	κ_2	λ_1	τ_0	M
$0 - 10^{-6}$	1.7×10^{-6}	2.8×10^{-6}	10^{-3}	0.9	4.3×10^{-6}	0.1	1.6×10^{-3}
		2.8×10^{-5}	10^{-2}		4.3×10^{-7}	10	5.3×10^{-3}
$0.05 - 0.9$	0.9	4×10^{-6}	0.95	0.9	3×10^{-6}	0.2	2×10^{-3}
		3.6×10^{-3}	0.1		3×10^{-7}	20	6×10^{-2}

Table 2. The numerical values of the parameters of the superpotential and the modulus vev τ_0 in the units $M_P = 1$. We have considered $B(\xi = 0) = 1$ and $A \sim 10^{-10}$.

play important roles in the inverse seesaw mechanism that we will demonstrate in the next section. We emphasize that f can't be zero in order to account for the correct dynamics of the waterfall fields.

3.3 Loop corrections

Here we discuss the effects from one-loop correction to the inflation potential, and two loop correction contributing to the gauge η -problem (or Dvali problem) mentioned in [37]. We will show that they don't have significant effect on the inflation potential.

- The Coleman-Weinberg one-loop corrections [36] are given by

$$V_{1\text{-loop}}(\Sigma) = \frac{1}{64\pi^2} \text{Str} \left[m^4 \left(\log \left(\frac{m^2}{Q^2} \right) - \frac{3}{2} \right) \right], \quad (3.11)$$

where $\text{Str} m(\Sigma)^4 = \sum_i (-1)^{2J_i} (2J_i + 1) m_i^4$ is the supertrace taken over all inflaton dependent mass eigenvalues of states with spin J_i . The stabilized fields during the inflation have $m_i \sim H$. But $H^2 M_P^2 \sim V_{\text{inf}} \sim A M_P^4$, therefore $H^2 \sim 10^{-10} M_P^2$ during the inflation. Accordingly, the 1-loop correction $V_{1\text{-loop}} \sim \frac{H^4}{64\pi^2} \sim \mathcal{O}(10^{-22}) M_P^4$, which is negligible compared to the tree level potential for any relevant renormalization scale.

- For the two loop Dvali problem, it may arise in the context of hybrid inflation if there is a gauge non-singlet flat directions (Σ in our case) as well as a singlet S coupled to the non-singlet waterfall fields. This results in large vacuum energy by the F-term F_S that breaks SUSY during the inflation. Therefore a significant two loop contribution to Σ effective mass is estimated as [24]

$$\delta m^2 \sim \frac{g_{B-L}^4}{(4\pi)^4} \frac{|W_S|^2}{m_\phi^2}, \quad (3.12)$$

where m_ϕ is the SUSY preserving mass of the waterfall fields and g_{B-L} is the $B - L$ gauge coupling. If $\delta m \sim H$, this implies $|\eta| \sim 1$ and the slow-roll conditions are violated causing an η -problem called the “gauge η -problem” or “Dvali Problem”. However we will show that the two loop contributions are suppressed compared to the tree level inflaton mass during inflation. Similar to the situation in [18], SUSY breaking scale M_{SB} during inflation gets contributions from $F_S = \kappa_1 \mu M_P$ and F_{S_i} . However the dominant contribution comes from F_{S_i} . According to the non-renormalization

theorem, corrections should be proportional to powers of $M_{SB} \sim \langle F_{S_i} \rangle$. Therefore the corrections arise from the following contributions [18, 24],

$$\delta m_\Sigma^2 \sim \frac{g_{B-L}^4}{(4\pi)^4} \frac{M_{SB}^4}{M_{Z'}^2}, \quad \frac{g_{B-L}^4}{(4\pi)^4} \frac{M_{SB}^4 H}{M_{Z'}^2}, \quad \frac{g_{B-L}^4}{(4\pi)^4} \frac{M_{SB}^4 H^2}{M_{Z'}^2}, \quad (3.13)$$

where $M_{Z'} \sim g_{B-L} \langle \Sigma \rangle$ is the $U(1)_{B-L}$ gauge boson. Assuming that during inflation, $\Sigma \sim 1M_P$, $g_{B-L} \sim 0.1$, $\mu \sim 10^{-5}M_P$, $\kappa_1 \leq 0.1$, $\kappa_2 \sim \mathcal{O}(1)$, $\lambda_1 \sim 10^{-5}$, $\tau_0 \sim 1M_P$, we find

$$\delta m_\Sigma^2 \sim 10^{-14} M_P^2, \quad 10^{-19} M_P^2, \quad 10^{-24} M_P^2, \quad (3.14)$$

which are negligible compared to the tree level mass squared of the inflaton during the inflation $m_\Sigma^2 \sim \mu^2 \sim 10^{-10} M_P^2$. Hence Dvali problem is absent in our scenario.

4 SUSY breaking, reheating and neutrino masses

In this section we investigate the link between the inflation sector and low energy physics sectors. We study the stringent constraints arising from reheating scenarios on the parameter space that is allowed by inflation observables as demonstrated in the previous section. Also, we discuss the neutrino masses which are one of the important motivations of our model. They are connected to the inflation sector via S , S_1 , S_2 . As we will see, SUSY breaking and R -symmetry will play essential rule in providing a tiny mass term for $S_1 S_2$, hence TeV inverse seesaw mechanism can be realized, with Y_ν not being small.

The assigned $B - L$ and R charges in table 1, forbid the terms SNN , $\phi_1 NN$ as well as higher order operators containing NN .³ Therefore the Majorana mass term $M_N NN$ is suppressed. On the other hand Z_3 symmetry with the assigned charges in table 1, prevents terms such as $\phi_1 S_2^2$ and $\phi_2 S_1^2$ hence any operator containing S_1^2 and S_2^2 is forbidden. The relevant terms in the superpotential to the neutrino masses which are allowed by symmetry, are given by

$$W_\nu = Y_\nu L H_u N + Y_S S N S_1 + \frac{\lambda_3}{M_P} S \phi_1 N S_2 + \left[\frac{\lambda_2}{M_P} S^2 + \kappa_2 \left(\frac{\phi_1 \phi_2}{M_P} - \mu \right) \right] S_1 S_2, \quad (4.1)$$

where the last term is common between the inflation and neutrino sector, hence a clear connection between the inflation and low energy physics of neutrino arises. SUSY breaking effects generate a tadpole term in the scalar potential $-2\kappa_1 M^2 m_{3/2} S + h.c.$, causing a shift in the vev of S, ϕ_1, ϕ_2 in terms of the gravitino mass $m_{3/2}$, as indicated in [38–45]

$$\langle S \rangle \simeq \frac{m_{3/2}}{\kappa_1}, \quad \langle |\phi_1| \rangle = \langle |\phi_2| \rangle \simeq M \left(1 - \frac{m_{3/2}^2}{\kappa_1^2 M^2} \right). \quad (4.2)$$

As advocated in the previous section, the term proportional to S^2 in the superpotential (4.1) had an important role in flattening the inflation potential, and here it appears the third and fourth roles of S in reheating and generating the tiny neutrino masses. Indeed a mass term $\mu_S S_1 S_2$ will be induced due to the shifts in the vevs of S and ϕ_i in (4.2).

³Here operators such as $(\phi_1 \phi_2)^n \phi_1 NN$ and $(\phi_1 \phi_2)^n N S_1$ are not allowed by R -symmetry in the UV sector as we indicated in section 2. On the other hand the operator $S^2 \phi_1 NN$ allowed by the symmetry will be suppressed by M_P^2 and has no effect in the neutrino mass matrix.

4.1 Constraints from reheating scenarios

At early time, after the inflation ended, the inflaton fields $\tilde{S}_1, \tilde{S}_2$ are massless. However, they acquire masses $m_{S_i} \ll M$, due to SUSY breaking. Therefore reheating the universe occurs via the decay of the heavy inflation sector fields S and ϕ_i to $N(\tilde{N})$ and $S_i(\tilde{S}_i)$, according to the interactions in the superpotential (4.1). Moreover $\tilde{S}_1, \tilde{S}_2$ can decay to the SM particles and contribute to the reheating according to their masses. In this respect the reheating temperature is given as [57–62]

$$T_R = \left(\frac{10}{g_s}\right)^{1/4} \left(\frac{2\Gamma M_P}{\pi c}\right)^{1/2}, \quad (4.3)$$

where Γ is the total decay width of the inflation sector fields, c is a constant of $\mathcal{O}(1)$, and g_s is the effective number of light degrees of freedom in the thermal bath at T_R .

In models of supergravity inflation, an upper bound on the reheating temperature arises due to the gravitino overproduction problem [46–55]. Gravitinos can be produced in the early universe thermally or non-thermally. If gravitinos are the lightest supersymmetric particle (LSP), then they are stable and can be a dark matter candidate. However, for high T_R , the energy density of gravitinos can be larger than the present total energy density of the universe hence overcloses the universe. Since the number density of gravitinos is proportional to T_R , this leads to a bound on the reheating temperature in terms of $m_{3/2}$. On the other hand if gravitinos are not stable, they may decay into the LSP (which can be neutralinos) before the Big Bang Nucleosynthesis (BBN), and hence the dark matter abundance imposes a constraint on gravitino production. Moreover, they can decay during and after the BBN. Constraints from BBN on both stable and unstable gravitinos were discussed in [54]. If $m_{3/2} \lesssim 10$ TeV, most of gravitinos decay after the start of BBN. Therefore, preserving the BBN predictions, which is compatible with the observed abundances of the light elements ^3He , ^4He , ^6Li and D, requires small abundance of gravitinos. This imposes an upper bound on T_R depending on the MSSM mass spectrum and $m_{3/2}$ as follows

$$10^6 \text{ GeV} \lesssim T_R \lesssim 2 \times 10^{10} \text{ GeV}, \quad \text{for} \quad 300 \text{ GeV} \lesssim m_{3/2} \lesssim 3 \times 10^3 \text{ GeV}. \quad (4.4)$$

For stable gravitinos, the MSSM-LSP becomes the next to LSP (NLSP) that decays to gravitinos. If their life time is long enough, this may affect the abundances of the light elements. Therefore BBN constrains $m_{3/2}$.

Nonetheless, the above derived upper bounds on reheating temperature can be relaxed to $\sim \mathcal{O}(10^{14})$ GeV as indicated in [56, 60], where models of higgs inflation usually predict such a high T_R . It was pointed out that the gravitino production problem can be addressed if gravitinos are heavy enough $m_{3/2} > \mathcal{O}(30 \text{ TeV})$ such that they decay before the BBN, or very light as in gauge mediation SUSY breaking models, $m_{3/2} \ll 1 \text{ GeV}$ where an upper bound on $T_R \lesssim \mathcal{O}(10^{14} \text{ GeV})$ is not problematic with the BBN observations.

Now we discuss the constraints from the reheating scenario imposed on the model parameters. We extract the interactions of S_i , S and ϕ_i leading to decay into fermions and

Figure 6. The allowed regions in the plane (λ_3, Y_S) , corresponding to constraints on reheating temperature.

scalars, that are not suppressed and contribute to T_R , as follows

$$\mathcal{L}_{S_i} = \left(Y_S Y_\nu^* \frac{m_{3/2}}{\kappa_1} S_1 + \lambda_3 Y_\nu^* \frac{M m_{3/2}}{\kappa_1 M_P} S_2 \right) \tilde{L}^* H_u^* + \text{h.c.} \quad (4.5)$$

$$\mathcal{L}_{\phi_i} = \kappa_2^2 \mu \sqrt{\frac{\mu}{M_P}} (\phi_1 + \phi_2) (\tilde{S}_1 \tilde{S}_1^* + \tilde{S}_2 \tilde{S}_2^*) + \kappa_2 \sqrt{\frac{\mu}{M_P}} (\phi_1 + \phi_2) \bar{S}_1^c S_2^c + \text{h.c.} \quad (4.6)$$

$$\begin{aligned} \mathcal{L}_S = S \left[\frac{|Y_S|^2 m_{3/2}}{\kappa_1} (\tilde{S}_2 \tilde{S}_2^* + \tilde{N} \tilde{N}^*) + \left(Y_S Y_\nu^* \tilde{S}_1 \tilde{L}^* H_u^* + \lambda_3 Y_\nu^* \sqrt{\frac{\mu}{M_P}} \tilde{S}_2 \tilde{L}^* H_u^* + \text{h.c.} \right) \right. \\ \left. + 2\kappa_1 \kappa_2 \mu \tilde{S}_1 \tilde{S}_2 \right] + \left[S \left(Y_S \bar{N}^c S_1^c + \lambda_3 \sqrt{\frac{\mu}{M_P}} \bar{N}^c S_2^c \right) + \text{h.c.} \right] \end{aligned} \quad (4.7)$$

In case of TeV scale gravitinos, $T_R \lesssim 2 \times 10^{10}$ GeV hence $\Gamma \lesssim 10^3$ GeV. Studying the branching ratios of different decay channels with respect to the total decay width, one finds that for large Y_S , the dominant decay channel is $S \rightarrow \bar{N}^c S_1^c$, hence a constraint on $Y_S < 10^{-4}$. In this case, the channel $S \rightarrow \bar{N}^c S_2^c$, $\phi_i \rightarrow \tilde{S}_i \tilde{S}_j^*$ become the dominant. The latter introduces the constraints $\lambda_3 < 10^{-3}$ and $\kappa_2 < 0.01$. As advocated above, these constraints can be relaxed if gravitinos were light $m_{3/2} \lesssim \mathcal{O}(0.1 \text{ GeV})$ or very heavy $m_{3/2} \gtrsim \mathcal{O}(100 \text{ TeV})$.

In figure 6, we show the allowed region under the curves, for the couplings Y_S and λ_3 , where we have fixed μ from the scalar perturbations observations and $Y_\nu \sim \mathcal{O}(1)$. The red curve corresponds to $T_R \sim 2 \times 10^{10}$ GeV with $\kappa_1 \sim 10^{-5}$, $\kappa_2 \sim 10^{-3}$ and $\tau_0 \sim 10^{-3} M_P$. The solid blue(purple) curves correspond to the relaxation $T_R \sim 10^{13}$ GeV, when gravitinos

are too heavy or too light, with the respective values $\kappa_1 \sim 10^{-2}(10^{-3})$, $\kappa_2 \sim 10^{-2}(10^{-1})$ and $\tau_0 \sim 1(10) M_P$. For the dashed we have changed $\tau_0 \sim 0.1(0.01) M_P$ respectively.

4.2 Neutrino masses

The Lagrangian of the neutrino masses can be derived from the superpotential (4.1), after electroweak symmetry breaking, as follows

$$\mathcal{L}_\nu = m_D \bar{\nu}_L N^c + M_{R_1} \bar{N}^c S_1^c + M_{R_2} \bar{N}^c S_2^c + \mu_S \bar{S}_1^c S_2^c + \text{h.c.}, \quad (4.8)$$

$$\begin{aligned} \text{with } m_D &= Y_\nu v \sin(\beta), & M_{R_1} &= \frac{Y_S m_{3/2}}{\kappa_1}, & M_{R_2} &= \frac{\lambda_3 m_{3/2} M}{\kappa_1 M_P}, \\ \mu_S &= \frac{\lambda_2}{M_P} \langle S \rangle^2 + \frac{\kappa_2}{M_P} [\langle \phi_1 \phi_2 \rangle - M^2] \\ &= \frac{m_{3/2}^2}{\kappa_1^2 M_P} \left[\lambda_2 - \kappa_2 + \frac{\kappa_2 m_{3/2}^2}{2\kappa_1^2 M^2} \right], \end{aligned} \quad (4.9)$$

where v is the electroweak vev and $\tan \beta = \frac{\langle H_u \rangle}{\langle H_d \rangle}$. The last term in μ_S is suppressed by the super-heavy scale M^2 and dominant effects come from the first two terms. The neutrino mass matrix, in the basis (ν_L, N, S_1, S_2) (for one flavor), is given by

$$M_\nu = \begin{pmatrix} 0 & m_D & 0 & 0 \\ m_D & 0 & M_{R_1} & M_{R_2} \\ 0 & M_{R_1} & 0 & \mu_S \\ 0 & M_{R_2} & \mu_S & 0 \end{pmatrix}. \quad (4.10)$$

For the mass hierarchy $\mu_S \ll m_D \ll M_{R_i}$, the eigenvalues consist of a light mass corresponding to the active neutrino, a relatively light one corresponding a sterile neutrino and two heavy states as follows

$$\begin{aligned} m_{\nu_l} &\simeq \frac{m_D^2 \mu_S}{2M_{R_1} M_{R_2}}, \\ m_{\nu_1} &\simeq \frac{\mu_S \left(-\frac{m_D^2}{2M_{R_2}} - 2M_{R_2} \right)}{M_{R_1}} \sim \mu_S, \\ m_{\nu_{2,3}} &\simeq \pm \left(M_{R_1} + \frac{m_D^2 + M_{R_2}^2}{2M_{R_1}} \right) + \frac{M_{R_2} \mu_S}{M_{R_1}} \sim \pm (M_{R_1} + M_{R_2}). \end{aligned} \quad (4.11)$$

In table 3, we list three points in the parameter space that give neutrino mass of order $\mathcal{O}(0.1)$ eV, where upper bound on the sum over the neutrino masses is given by $\sum_\gamma m_{\nu_\gamma} < 0.15$ eV, at 95% CL [63, 64]. We have taken different value of $m_{3/2}$ corresponding to different regimes in the reheating scenario and fixed $\sin(\beta) \sim 1$. The value of M can range between $10^{15} - 10^{17}$ GeV. Figure 7 depicts the allowed values of the parameters in the plane (λ_3, Y_S) , while fixing the other parameters, where the solid curves correspond to active neutrino mass value $|m_{\nu_l}| \sim \mathcal{O}(0.1)$ eV. We have fixed $\sin(\beta) \sim 1$, and $Y_\nu \sim \mathcal{O}(1)$ in the case where $m_{3/2} = 0.1, 10^5$ GeV. We have imposed the constraints from inflation observables and reheating temperature according to the cases discussed in (4.1). Therefore the segments between the black dots give the allowed values in the parameter space (λ_3, Y_S) for fixed values of the other parameters.

Figure 7. The regions of parameter space in the plane (λ_3, Y_S) allowed by neutrino masses and reheating temperature constraints. The solid lines represent the lightest neutrino eigenstate $|m_{\nu_1}| \sim \mathcal{O}(0.1)$ eV, while the dotted curve represent the constraints from reheating.

$m_{3/2}$	λ_2	λ_3	κ_1	κ_2	Y_ν	Y_S	$ m_{\nu_1} $	$ m_{\nu_{2,3}} $
10^3	0.008	0.002	10^{-5}	10^{-3}	0.05	3×10^{-5}	3×10^{-5}	4.4×10^3
0.1	0.9	0.4	10^{-7}	0.1	0.9	0.01	2.1×10^{-7}	1.06×10^4
10^5	0.1	0.01	0.01	0.01	0.5	0.0032	6.6×10^{-7}	3.21×10^4

Table 3. Three benchmark points in the parameter space that give the active neutrino masses $|m_{\nu_i}| \sim \mathcal{O}(0.1)$ eV and satisfy the constraints from inflation observables and reheating. All mass scales are given in GeV.

5 Conclusions

We have constructed an inflation model in which the non-singlet right-handed sneutrinos S_1, S_2 contain the inflaton component, in no-scale supergravity. We have modified the tribrid inflation model to implement a scenario in which the inflation sector is linked to a TeV inverse seesaw mechanism. In this regard, we have explained the smallness of the μ_S mass parameter on a basis of the symmetry defined on the model and SUSY breaking effects. We have discussed different regimes in the parameter space that lead to successful inflation compatible with Planck observation. We found that the reheating scenario imposes severe constraints on the parameter space when gravitino mass is TeV scale. These constraints can be relaxed for very heavy or ultra-light gravitinos in order to avoid the BBN constraints. We have derived the neutrino masses taking into account the latter constraints. The model may have interesting consequences on the low energy phenomenology.

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