Sub-Planckian Inflation Due To A Complex Scalar In A Modified Radiative Seesaw Model

DISSERTATION ABSTRACT

Sub-Planckian Inflation Due To A Complex Scalar In A Modified Radiative Seesaw Model

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I Research Motivation

The standard model (SM) of particle is now considered to be extended due to unsuccessful explanation of some observational phenomena in this framework. Those phenomena are the neutrino masses and mixing [1], the existence of dark matter [2] and the baryon number asymmetry in the universe [3]. Finding a model that can explain all those phenomena simultaneously without causing any tension to other phenomenological problems such as lepton flavor violating processes would be a crucial step to understand the new physics beyond the SM. One of a promising candidate for that purpose is a simple extension of the SM with an inert doublet scalar and three right-handed neutrinos. Several studies [4–6] show that possibility.

On the other hand, the existence of inflationary expansion of the universe at very early time is strongly supported by the CMB observations. Severe observational constraints such as Planck 2013 and Bicep2 restrict the allowed inflation model now [7, 8]. They completely disfavor any model predicting at almost scale invariant and blue tilted scalar power spectrum. They also prefer to a single field model over more complicated scenarios. There are also theoretical constraints such as the Lyth bound [9] that restricts the allowed field value to realize the sufficient tensor-to-scalar ratio. The η problem is another one that is a kind of hierarchy problem between the inflaton mass and the Hubble parameter. In single field inflation models, since the Lyth bound prevents the inflaton field to have a value below Planck scale, the higher order terms suppressed by the Planck mass appear to ruin the flatness of the inflaton potential. If there is no symmetry protecting the potential, this difficulty is caused and the η problem is inevitable as well. The observation by Planck 2015 [10] tightens the tensor-to-scalar ratio constraint to be $r_{0.002} < 0.11$ (95 % CL) so that only a few model can still survive, as instances the hiltop quartic model, R^2-inflation, Higgs-inflation and power-law chaotic inflation with power less than two.

From such many inflation models that survive from the observational constraints, there are not so many inflaton candidates that play any role in particle physics. Even so, they have still problems. As instances, the power-law chaotic inflation which is motivated by axion monodromy suffers trans-Planckian problem due to the Lyth bound and the η problem, and the Higgs inflation suffers from the unitary problem caused by a large non-minimally coupling [11, 12].

Motivated by the above facts, we consider an extension of the radiative seesaw model with a complex scalar to explain the inflation of the universe as well without disturbing favorable features of the original model. To evade the Lyth bound and the η problem, the field value of the inflaton which corresponds to the complex scalar will be kept in sub-Planckian values by choosing a potential in such a way that only a particular dynamics of the inflaton is allowed. In this scenario, the spectral index and the tensor-to-scalar ratio could have values in a region favorable by the recent CMB observations depending on the parameter sets in the inflaton potential.

II Modification of The Radiative Neutrino Mass Generation Model

Radiative seesaw scenario is an alternative way to explain tiny neutrino masses. In this scenario they are radiatively induced at the one loop level by imposing an exact Z_2 symmetry and introducing additional Z_2-odd scalar doublet η and Z_2-odd right-handed neutrino $N_i (i = 1, 2, 3)$ [13]. All of the standard model particle are labeled by even parity. As a result of this assignment, the Lagrangian of this model is

$$-\mathcal{L}_N = -h_{\alpha i}\bar{N}_i\eta^\dagger l_\alpha - h^*_{\alpha i}\bar{l}_\alpha \eta N_i + \frac{M_i}{2}\bar{N}_i N_i^c + \frac{M_i^*}{2}\bar{N}_i^c N_i \tag{1}$$

with the scalar sector potential is given as

$$\begin{aligned} V_{\text{scalar}} = \; & m_1^2\Phi^\dagger\Phi + m_2^2\eta^\dagger\eta + \lambda_1(\Phi^\dagger\Phi)^2 + \lambda_2(\eta^\dagger\eta)^2 + \lambda_3(\Phi^\dagger\Phi)(\eta^\dagger\eta) + \lambda_4(\Phi^\dagger\eta)(\eta^\dagger\Phi) \\ & + \frac{1}{2}\left[\lambda_5(\Phi^\dagger\eta)^2 + \lambda_5^*(\Phi\eta^\dagger)^2\right]. \end{aligned} \tag{2}$$

Any bilinear term $(\Phi^\dagger\eta)$ is forbidden by the Z_2 symmetry so that λ_5 can always be chosen as a real parameter by the field redefinition for η. Under the assumption that $m_1^2 < 0$ and $m_2^2 > 0$, Higgs Φ obtains the vacuum expectation value $v := \sqrt{-m_1^2/2\lambda_1} = \langle\phi_0\rangle$ [15].

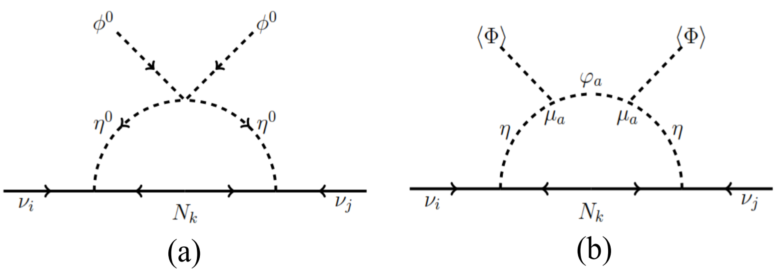


FIGURE 1: (a) One-loop generation of neutrino mass considered in the radiative neutrino masses model with an inert doublet [14] (b) One-loop generation of neutrino masses in the present model. The coupling μ_a in this diagram is defined as $\mu_1 := \frac{\mu}{\sqrt{2}}$ and $\mu_2 := \frac{i\mu}{\sqrt{2}}$

Neutrino mass is generated through the one loop diagram given by Fig 1.a, that involves the exchange of η_R^0 and η_I^0. Applying Feynman rules to the diagram gives the neutrino mass matrix :

$$\mathcal{M}_{ij}^{\nu} = \frac{h_{ik}h_{kj}}{16\pi^2} M_k \left[\frac{m_{\eta_R^0}^2}{m_{\eta_R^0}^2 - M_k^2} \ln\left(\frac{m_{\eta_R^0}^2}{M_k^2}\right) - \frac{m_{\eta_I^0}^2}{m_{\eta_I^0}^2 - M_k^2} \ln\left(\frac{m_{\eta_I^0}^2}{M_k^2}\right) \right]. \tag{3}$$

where $m_{\eta_{R,I}^0}$ denote the mass of the neutral components of the inert doublet and M_k denotes the mass of the right-handed neutrinos. If we use the quantities $\Delta m^2 := (m_{\eta_R^0}^2 - m_{\eta_I^0}^2)/2 = \lambda_5 v^2$ and $m_0^2 := (m_{\eta_R^0}^2 + m_{\eta_I^0}^2)/2$, under the assumptions that $m_0^2 \gg \Delta m^2$, $m_{\eta_R^0}^2 \simeq m_{\eta_I^0}^2 \simeq m_0^2$ are satisfied, the neutrino mass is approximated as

$$\mathcal{M}_{ij}^{\nu} = \sum_{k=1}^{3} \frac{h_{ik}h_{kj}M_k}{8\pi^2} \frac{\lambda_5 v^2}{(m_0^2 - M_k^2)} \left[1 - \frac{M_k^2}{(m_0^2 - M_k^2)} \ln\left(\frac{m_0^2}{M_k^2}\right) \right]. \tag{4}$$

This equation shows that the smallness of λ_5 is a crucial role to explain the smallness of neutrino masses for the TeV range M_k and m_0.

To explain the smallness of λ_5 in the radiative seesaw model, we can consider a scenario in which this coupling is an effective coupling in low energy region resulted from integrating out of a heavy singlet scalar S. In this scenario, the coupling λ_5 in the original model is supposed to be zero. We will explain how λ_5 is derived from the extended model later. The new singlet scalar S should be a Z_2 odd field in order to couple with the inert doublet scalar η and Higgs doublet scalar Φ. The additional Lagrangian terms should be added in the original model are

$$\begin{aligned} -\mathcal{L}_S = \quad & \tilde{m}_S^2 S^\dagger S + \frac{1}{2} m_S^2 S^2 + \frac{1}{2} m_S^2 S^{\dagger 2} + \kappa_1 \left(S^\dagger S\right)^2 + \kappa_2 \left(S^\dagger S\right)\left(\Phi^\dagger \Phi\right) + \kappa_3 \left(S^\dagger S\right)\left(\eta^\dagger \eta\right) \\ & - \mu S \eta^\dagger \Phi - \mu^* S^\dagger \Phi^\dagger \eta. \end{aligned} \tag{5}$$

Writing $S = \frac{1}{\sqrt{2}}(\varphi_1 + i\varphi_2)$ the masses of each components are given as $\bar{m}_1^2 = \tilde{m}_S + m_S^2$ and $\bar{m}_2^2 = \tilde{m}_S^2 - m_S^2$. Since Z_2 is considered to be an exact symmetry, $\tilde{m}_S^2 > m_S^2$ is satisfied.

Neutrinos still remain massless at tree level such as in the original Ma-model since η has zero vacuum expectation value. However, neutrino masses can be generated through diagram given in the Figure 1.b. By applying Feynman rules, the resulting neutrino mass matrix is found as

$$(\mathcal{M})_{\alpha\beta} = \sum_{i=1}^{3} \sum_{a=1}^{2} \frac{h_{\alpha_i} h_{\beta i} \mu_a^2 \langle \Phi \rangle^2}{8\pi^2} I(M_i^2, M_\eta^2, \bar{m}_a^2). \tag{6}$$

where $I(M_i^2, m_1^2, m_2^2)$ is defined as

$$I(M_i^2, m_1^2, m_2^2) = -\frac{M_i^2 \ln(M_i^2)}{(M_i^2 - m_2^2)(M_i^2 - m_1^2)^2} + \frac{(m_1^4 - M_i^2 m_2^2)\ln(m_1^2)}{(M_i^2 - m_1^2)^2(M_i^2 - m_2^2)^2} + \frac{m_2^2 \ln(m_2^2)}{(m_1^2 - m_2^2)^2(M_i^2 - m_2^2)} - \frac{1}{(M_i^2 - m_1^2)(m_1^2 - m_2^2)}. \tag{7}$$

This masses matrix is reduced to the neutrino masses matrix of Ma-model under assumption that the condition $\tilde{m}_S \gg m_S, m_\eta, M_i$ is satisfied. The approximated formula is given by

$$(\mathcal{M})_{\alpha\beta} = \sum_{i=1}^{3} \frac{h_{\alpha_i} h_{\beta i} \langle \Phi \rangle^2}{8\pi^2} \frac{m_S^2 \mu^2}{\tilde{m}_S^4} \frac{M_i}{M_\eta^2 - M_i^2} \left[1 + \frac{M_i^2}{M_\eta^2 - M_i^2} \ln\left(\frac{M_i^2}{M_\eta^2}\right) \right], \tag{8}$$

where the factor $\frac{m_S^2 \mu^2}{\tilde{m}_S^4}$ appears from $\sum_{a=1}^{2} \mu_a^2 / \bar{m}^2$. Comparing this to equation (4), it is obvious that the coupling constant λ_5 for the $(\eta^\dagger \Phi)^2$ in the original model is effectively approximated as the quantity $\frac{m_S^2 \mu^2}{\tilde{m}_S^4}$. We might interpret the original model as the low energy limit of the present extended model, in which λ_5 is an effective coupling derived from the interaction $-\mu S \eta^\dagger \Phi - \mu^* S^\dagger \Phi^\dagger \eta$ by integrating out S. At tree level of this extended model, the amplitude of the interaction $\eta\Phi \to \eta\Phi$ is given by

$$\mathcal{M} \simeq \left[\frac{\mu_1^2}{(q^2 - \bar{m}_1^2)} - \frac{\mu_2^2}{(q^2 - \bar{m}_2^2)} \right] \simeq \mu^2 \left[\frac{m_S^2}{\bar{m}_1^2 \bar{m}_2^2} \right]_{\bar{m}_1^2, \bar{m}_2^2 \gg q^2} \simeq \frac{\mu^2 m_S^2}{\tilde{m}_S^4}, \tag{9}$$

which coincides with λ_5 in the original Ma-model. Hierarchical masses problem bet–ween μ, m_S and $\tilde{m}_S$ now replaces the smallness problem of λ_5 in the Ma-model. It is a key factor to explain the smallness of the neutrino masses. If we leave the origin of this hierarchy problem to a complete theory at high energy regions, all the neutrino masses, the DM abundance and the baryon number asymmetry could be also explained in this extended model at TeV regions just as discussion given in [5].

III Aspects as The Inflation Model

III.1 General features of the model

We consider an inflation scenario working at sub-Planckian scale by introducing non-renormalizable terms obeying Z_2 symmetry to the potential for complex scalar field S given in equation (5). These terms could restrict the trajectory of the evolution of S. In that case, even though the radial motion of S is small, additional angular motion makes its whole trajectory length sufficiently large to evade the Lyth bound.

As such an example, lets assume that the complex scalar S has Z_2 invariant additional potential as below

$$V = c_1 \frac{\left(S^\dagger S\right)^n}{M_{\rm pl}^{2n-4}} \left[1 + c_2 \left(\frac{S}{M_{\rm pl}}\right)^{2m} \exp\left(i\frac{S^\dagger S}{\Lambda^2}\right) + c_2 \left(\frac{S^\dagger}{M_{\rm pl}}\right)^{2m} \exp\left(i\frac{S^\dagger S}{\Lambda^2}\right)\right] \tag{10}$$

$$= c_1 \frac{\varphi^{2n}}{2^n M_{\rm pl}^{2n-4}} \left[1 + 2c_2 \left(\frac{\varphi}{\sqrt{2} M_{\rm pl}}\right)^{2m} \cos\left(\frac{\varphi^2}{2\Lambda^2} + 2m\theta\right)\right], \tag{11}$$

where $M_{\rm pl}$ is the reduced Planck mass, and both of n and m are positive integers. In the second line, we adopt polar coordinate expression for $S = \frac{1}{\sqrt{2}}\varphi e^{i\theta}$. The most crucial part in the potential is the exponential term. However, we cannot explain its origin in this stage. We only expect that it might be effectively induced through the nonperturbative dynamics in the UV completion of the model.

The quantities characterizing the inflation need to be calculated to understand the features of the inflation. Some of them are the slow-roll parameters η and ε, the e-folding number N, the spectral index n_s, the tensor-to-scalar ratio r and the running of the spectral index n'_s. If the inflaton is restricted to move along the minimums of the potential, The condition $\Lambda \ll \varphi < M_{\rm pl}$ is satisfied. As results, those mentioned quantities are given as follow:

$$\varepsilon := \left(M_{\rm pl}^2/2\right)\left(V'/V\right)^2 = m^2 \bar{\Lambda}^4 \bar{\varphi}^{-6} \left[B(\varphi)/A(\varphi)\right]^2, \tag{12}$$

$$\eta := M_{\rm pl}^2 \left(V''/V\right) = m^2 \bar{\Lambda}^4 \bar{\varphi}^{-6} \left[C(\varphi)/A(\varphi)\right], \tag{13}$$

$$\xi := M_{\rm pl}^4 \left(V'V'''/V^2\right) = m^4 \bar{\Lambda}^8 \bar{\varphi}^{-12} \left[B(\varphi)D(\varphi)/A(\varphi)^2\right] \tag{14}$$

$$n_s \simeq 1 - 6\varepsilon + 2\eta, \quad r \simeq 16\varepsilon, \quad n'_s := \frac{dn}{d\ln k} \simeq 16\varepsilon\eta - 24\varepsilon^2 - 2\xi \tag{15}$$

$$N := N(\varphi) - N(\varphi_e) \tag{16}$$

where we have defined $\bar{\varphi} := \left(\frac{\varphi}{\sqrt{2}M_{\rm pl}}\right)$, $\bar{\Lambda} := \left(\frac{\Lambda}{M_{\rm pl}}\right)$ and

$$A(\varphi) := 1 - 2c_2\bar{\varphi}^{2m}, \tag{17}$$

$$B(\varphi) := n - 2c_2(n+m)\bar{\varphi}^{2m}, \tag{18}$$

$$C(\varphi) := n(2n-3) - 2c_2(n+m)(2n+2m-3)\bar{\varphi}^{2m}, \tag{19}$$

$$D(\varphi) := n(2n-3)(2n-6) - 2c_2(n+m)(2n+2m-3)(2n+2m-6)\bar{\varphi}^{2m}, \tag{20}$$

$$N(\varphi) := \frac{\bar{\Lambda}^{-4}\bar{\varphi}^6}{6nm^2}\left[1 + \frac{6c_2 m\bar{\varphi}^{2m}}{n(3+m)} \; F\left(1, \frac{3}{m}+1; \frac{3}{m}+2; 2c_2\left(1+\frac{m}{n}\right)\bar{\varphi}^{2m}\right)\right]. \tag{21}$$

The quantity $F(a,b;c;x)$ denotes the Hypergeometric function. The Energy scale of the inflation is fixed by the normalization of scalar power spectrum $\Delta_{\mathcal{R}}^2 := \frac{V}{24\pi^2 M_{\rm pl}^4 \varepsilon}\Big|_{k^*}$, which is given about $\ln(10^{10}\Delta_{\mathcal{R}}^2) = 3.094 \pm 0.034$ (68% CL, *Planck* TT, TE, EE+lowP combination data) [10]. All of the slow-roll parameters and cosmological parameters compared to observations should be represented at the horizon exit $k_* = aH$.

To understand whole processes that the scalar field S will undergo, it is useful to investigate the time evolution of the fields numerically. Each component field $\varphi_{1,2}$ of the scalar field $S = \frac{1}{\sqrt{2}}(\varphi_1 + i\varphi_2)$ follows the equation of motions as follows:

$$\ddot{\varphi}_i + 3H\dot{\varphi}_i = -\frac{\partial V}{\partial \varphi_i} \quad (i = 1, 2), \tag{22}$$

where the Hubble parameter H of the system is now written as $H^2 = \frac{1}{3M_{\rm pl}^2}\left(\sum_i \frac{1}{2}\dot{\varphi}_i + V\right)$ and $\frac{\partial V}{\partial \varphi_i}$ denotes the partial derivative of potential $V(S)$ in the direction of the field component φ_i. The initial value of the field could not be selected arbitrarily as it could ruin its dynamics depending on it. The best way is to place the inflaton initial value at potential minimum. At a particular point, the motion suddenly falls toward of the center of the potential and starts the oscillation. This point is considered to the time when the inflation end and the reheating after inflation takes place to convert energy density of the inflaton to the particles production. It is related to the time when slow-roll parameter $\varepsilon(t) := -\dot{H}/H^2$ is close to unity but mostly much less than unity.

III.2 Constraints from Planck 2013, Bicep2 and Planck 2015

The Planck 2013 constrains the scalar spectral index to be $n_s = 0.9603 \pm 0.0073$ and the scalar power spectrum amplitude to be $\Delta_{\mathcal{R}}^2 = 2.196^{+0.051}_{-0.06} \times 10^{-9}$. It also and establishes an upper bound on the tensor-to-scalar ratio as $r < 0.11$ (95% CL) at the pivot scale $k_* = 0.002$ Mpc^{-1}. Thus, any models predicting subtantial deviation from the nearly scale invariance and a blue tilted scalar power spectrum, such as a original hybrid model, are ruled out. Moreover, since the level of local non-Gaussianities is constrained by a bound $f_{\rm NL}^{\rm loc} = 2.7 \pm 5.8$, Planck 2013 data prefers single field inflation to more complicated possibilities. Some of them, i.e. exponential potential models, the simplest hybrid inflationary models, and monomial potential models of degree $n \geq 2$, do not provide a good fit to the data [7]. However, Bicep2 measures rather higher tensor-to-scalar ratio than that measured by Planck 2013. The bound is given at $r = 0.20^{+0.07}_{-0.05}$ without dust foreground subtraction that disfavors $r = 0$ at 7.0σ level, or it is given at $r = 0.16^{+0.06}_{-0.05}$ if dust foreground subtraction is included [16]. Monomial potential with power $2 \leq p \leq 3$ survives under this new constrain. Here we give more attention to $p = 2$ which corresponds to $n = 3$ in our model to minimize the tension between the results of Planck 2013 and Bicep2. Predictions given for some parameter sets for $n = 3$ and $m = 1$ are presented in the Table 1 and in the Figure 2.

	c_1	c_2	$\frac{\Lambda}{M_{\rm pl}}$	$\frac{\varphi^*}{\sqrt{2}M_{\rm pl}}$	H^*(GeV) $\times 10^{14}$	N^*	n_s	r	n'_s
A	1.66×10^{-6}	0.7	0.04	0.378	0.871	59.0	0.971	0.107	-0.00016
	2.04×10^{-6}	0.7	0.04	0.371	0.921	54.2	0.968	0.119	-0.00022
	2.42×10^{-6}	0.7	0.04	0.366	0.965	49.1	0.965	0.131	-0.00027
B	0.257	6.0	0.002	0.0512	0.945	60.4	0.969	0.124	-0.00046
	0.305	6.0	0.002	0.0505	0.986	55.0	0.966	0.136	-0.00054
	0.364	6.0	0.002	0.0498	1.030	50.0	0.962	0.149	-0.00064

TABLE 1: Predictions for some typical parameter sets of the model defined for $n = 3$ and $m = 1$.

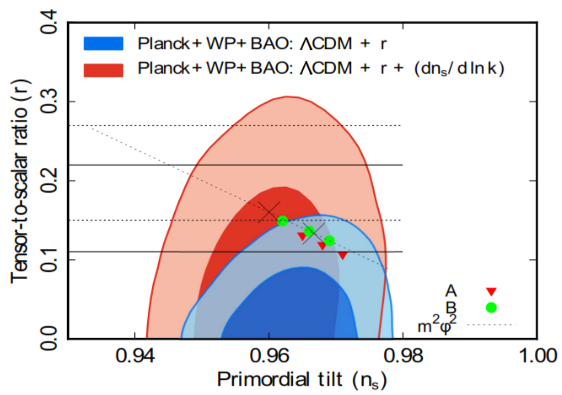
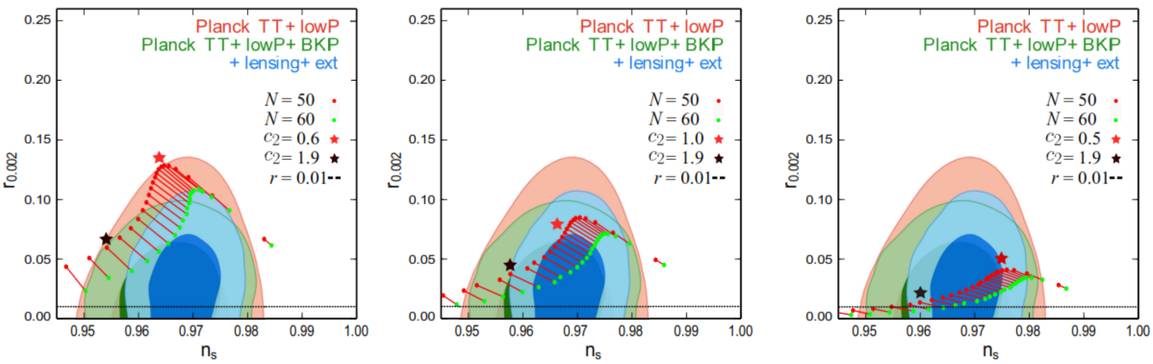


FIGURE 2: Predicted values of (n_s, r) for several parameter sets $\left(c_2, \frac{\Lambda}{M_{\rm pl}}\right)$ given in the Table 1 are plotted here. The dotted line represents the prediction of the quadratic chaotic inflation model, in which the points corresponding to $N_* = 50$ and 60 are represented as crossed lines. The horizontal solid lines and dotted lines represent the Bicep2 1σ constraints with and without the foreground subtraction, respectively [16]. The contours given as Figure 4 in Planck Collaboration XXII [7] are used here. Since the running of the spectral index is negligible, the blue contour should be compared with the predictions

The Planck 2015 mission releases the announcement that the spectral index of curvature perturbations is measured to be $n_s = 0.968 \pm 0.006$ with the tight constraint of scale dependence $dn_s/d\ln k = 0.003 \pm 0.007$. The upper bound on the tensor-to-scalar ratio is $r < 0.11(95\%\ \text{CL})$ measured at pivot scale $k_* = 0.002\ \text{Mpc}^{-1}$ which is stronger than before. Even so, monomial inflation with power $p < 2$ is found to survive from the constraint [10]. Due to this finding, we need to consider new parameter sets given for $n \leq 3$ to find new predictions in this model. The complete predictions are plotted in the Figure 3. This figure shows that $n = 1$ is the most favorable one in the present model.



(a) (b) (c)

FIGURE 3: Predicted regions in the (n_s, r) plane are presented in panel (a) for $n = 3$, in panel (b) for $n = 2$, and in panel (c) for $n = 1$. Λ is fixed as $\Lambda = 0.05 M_{\rm pl}$ in all cases. Contours given in the right panel of Fig. 21 in Planck 2015 results.XIII.[17] are used here. Horizontal black lines $r = 0.01$ represent a possible limit detected by LiteBIRD in near future.

III.3 Reheating after inflation

The early stage of the reheating may be constituted by two main processes: preheating due to the parametric resonance through quartic interactions of S with Φ and η and the perturbative decay due to an interaction term $\mu S \eta^\dagger \Phi$. The first constituent may not occur effectively as the fields coupled to $\varphi_{1,2}$ have large effective mass so it seems difficult for $\varphi_{1,2}$ to produce these particles. The perturbative decay due to $\frac{\mu}{\sqrt{2}}\varphi_1 \eta^\dagger \Phi$ and $\frac{i\mu}{\sqrt{2}}\varphi_2 \eta^\dagger \Phi$ takes place to complete energy transfer from the inflaton to the radiation. The decay width of each process is given by $\Gamma_{\varphi_i} = \frac{1}{8\pi}\frac{|\mu|^2}{\bar{m}_i}$ where $\bar{m}_1^2 = \tilde{m}_S^2 + m_S^2$ and $\bar{m}_2^2 = \tilde{m}_S^2 - m_S^2$ are the mass of φ_1 and φ_2, respectively. As $\tilde{m}_S$ is assumed much larger than m_S, the reheating temperature given from this perturbative decay can be estimated as [18]

$$T_R \simeq 1.6 \times 10^8 \left(\frac{|\lambda_5|}{10^{-6}}\right)^{1/2} \left(\frac{\tilde{m}_S}{m_S}\right) \sqrt{\frac{\tilde{m}_S}{10^6 \text{ GeV}}} \quad \text{GeV}. \tag{23}$$

If we taking the lightest neutral component of η as dark matter with mass of order 1 TeV, it suggests that $|\lambda_5|$ should be $O(10^{-6})$ or less [20, 21]. Thus the reheating temperature would be vary in the range of 10^5 GeV $\leq T_R \leq 10^{15}$ GeV depending on the value of $\tilde{m}_S$. This order is high enough to produce termal right-handed neutrinos of $O(1)$ TeV to produce sufficient baryon asymmetry via leptogenesis.

IV Conclusion

An extension of the radiative neutrino masses model by a complex singlet has been considered to explain the inflation of the universe by keeping favorable features of the original model, the simultaneous explanation of the small neutrino masses, the DM abundance and the baryon number asymmetry in the Universe. The complex singlet not only plays a role in the inflation scenario due to its component but is also involved in the neutrino mass generation at one loop to explain the smallness of neutrino masses. By choosing a complex scalar potential realizing a dynamics of the inflaton following a spiral-like valley, trans-Planckian field variation can be realized to generate the sufficient e-foldings even though the relevant field is kept sub-Planckian. The η problem is now stated in the different way, that is the mass hierarchy of $\tilde{m}_S^2, m_S^2, \kappa\varphi^2 \ll H^2$ which is relevant to the neutrino mass and the scale hierarchy $\Lambda \ll M_{\rm pl}$. The UV completion of the model is expected to give a solution for it. The origin of the potential cannot be still discovered at this stage.

The model interestingly behaves like a single field inflation scenario which is closely related with the power-low chaotic inflation in a limiting case. We have shown that the predicted values for them by using the parameter sets for $n = 1, 2$ and $m = 1$ are favorable even for Planck 2015 observational constraints.

Furthermore, the rough estimation of the reheating temperature in this model could be high enough to produce thermal right-handed neutrinos for resonant leptogenesis. Therefore, the model seems to have no serious difficulty to explain the crucial problems beyond the SM including the baryon number asymmetry like the original model of the radiative neutrino masses model.

References

[1] Fukuda, Y. et al. [Super-Kamiokande Collaboration], Phys.Rev.Lett. 81 (1998) 1562-1567.

[2] Tegmark, Max et al. [SDSS Collaboration], Phys.Rev. D69 (2004) 103501.

[3] Riotto, Antonio et al, Ann.Rev.Nucl.Part.Sci. 49 (1999) 35-75 .

[4] T. Hambye, F.-S. Ling, L. Lopez Honorez, and J. Rocher, JHEP 1005 (2010) 066, arXiv:0903.4010 [hep-ph].

[5] Daijiro Suematsu, Phys.Rev. D85 (2012) 073008.

[6] Shoichi Kashiwase and Daijiro Suematsu, Eur.Phys.J. C73 (2013) 2484.

[7] P. Ade et al., arXiv:1303.5082 [astro-ph.CO], 2013.

[8] P. A. R. Ade et al. [BICEP2 Collaboration], arXiv:1403.4302 [astro-ph.CO].

[9] David H. Lyth, Phys.Rev.Lett. 78 (1997) 1861-1863.

[10] P. A. R. Ade et al., arXiv:1502.02114v1 [astro-ph.CO].

[11] Lerner, Rose N. et al, Phys.Rev. D83 (2011) 123522.

[12] Lerner, Rose N. et al., JCAP 1211 (2012) 019.

[13] S. Kanemura, T. Matsui and T. Nabeshima, Phys. Lett. B 723 (2013) 126.

[14] E. Ma, Phys. Rev. D 73, 077301 (2006).

[15] Ethan M. Dolle and Shufang Su, Phys.Rev. D80 (2009) 055012 .

[16] P. A. R. Ade et al. [BICEP2 Collaboration], arXiv:1403.3985 [astro-ph.CO].

[17] (P.A.R. Ade (Cardiff U.) et al, arXiv:1502.01589 [astro-ph.CO].

[18] Lev Kofman, Andrei Linde and Alexei A. Starobinsky, arXiv:hep-ph/9704452v2.

[19] Rouzbeh Allahverdi, Robert Brandenberger, Francis-Yan Cyr-Racine and Anupam Mazumdar, arXiv:1001.2600 [hep-th].

[20] S. Kashiwase and D. Suematsu, Phys. Rev. D86 (2012) 053001.

[21] S. Kashiwase and D. Suematsu, Eur Phys. J C73 (2013) 2484.

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学位論文審査報告書（甲）

1．学位論文題目（外国語の場合は和訳を付けること。）

Sub-Planckian Inflation Due To A Complex Scalar In A Modified Radiative Seesaw Model

（修正輻射シーソーモデルでの複素スカラーによるサブプランクスケールインフレーション）

2．論文提出者　(1) 所　　属　　数物科学　　　　専攻

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3．審査結果の要旨（600～650 字）

近年の宇宙背景放射の詳細な観測により、宇宙初期にインフレーションと呼ばれる指数関数的な宇宙膨張の存在が示唆されている。これまで多数のインフレーション模型が提案されてきたが、これらの観測に基づきその多くは既に棄却される一方、現実的な模型は未だ構成されていない。

本論文では、ニュートリノ質量生成機構を持つ模型として提案された輻射シーソー模型を拡張することで、ニュートリノ質量生成と密接な関わりを持つスカラー場がインフラトンの役割を担うインフレーション模型を提案している。この模型において以下の結果が示されている。(1) プランク質量より小さなインフラトンの値で十分なインフレーションを実現できることを示すことで、Lyth 限界問題への一つの解を与えた。(2) 宇宙背景放射の揺らぎに関するスカラー指標やテンソル・スカラー比に観測結果と整合する予言値を与えた。(3) インフラトンにニュートリノ質量生成という素粒子論的な役割を担わせることで、再加熱以降の物理の解析を可能とした。これらの結果は、素粒子模型とインフレーションを関連付ける興味深い可能性を提示しており、今後のインフレーション模型の研究に有用な知見をもたらすものと考えられ、博士論文に値すると判断する。

4．審査結果　(1) 判　　定（いずれかに○印）　合　格 ・ 不合格

(2) 授与学位　　博　士（　理学　）