

Dissertation Abstract

***Inflation and Reheating With The Singlet Scalars Related to
CP Violation***

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Abstract

We proposed a simple extension of the standard model which contains new fermions and scalars. This extension will perhaps solve the problems of the early universe and also explain the origin of the CP violation. In the inflation cosmology, we proposed that inflaton is the linear combination of the introduced singlet scalars. This new-defined inflaton gives the new features in both inflation and reheating scenarios. Also, we expect the new fermions can give the suitable parameters which solves the phase in CKM and PMNS matrices. Finally, we show that both introduced fermions and scalars can have a new feature in low-scale leptogenesis.

keywords singlet scalars, inflation, leptogenesis, CP violation

I. THE LAGRANGIAN

We will start our discussion with the introduction of the Lagrangian

$$\begin{aligned}
-\mathcal{L}_Y = & y_D \sigma \bar{D}_L \bar{D}_R + y_E \sigma \bar{E}_L \bar{E}_R + \sum_{j=1}^3 \left[\frac{y_{N_j}}{2} \sigma \bar{N}_j^c N_j + y_{D_j} S \bar{D}_L d_{R_j} + \tilde{y}_{d_j} S^\dagger \bar{D}_L d_{R_j} \right. \\
& \left. + y_{e_j} S \bar{E}_L e_{R_j} + \tilde{y}_{e_j} S^\dagger \bar{E}_L e_{R_j} + \sum_{\alpha=1}^3 h_{\alpha j}^* \eta \bar{l}_\alpha N_j \right]
\end{aligned} \tag{1}$$

and potentials

$$\begin{aligned}
V = & \lambda_1 (\mathcal{H}^\dagger \mathcal{H})^2 + \lambda_2 (\eta^\dagger \eta)^2 + \lambda_3 (\mathcal{H}^\dagger \mathcal{H})(\eta^\dagger \eta) + \lambda_4 (\mathcal{H}^\dagger \eta)(\eta^\dagger \mathcal{H}) + \frac{\lambda_5}{2M_*} [\sigma(\eta^\dagger \mathcal{H})^2 + \text{h.c.}] \\
& + \kappa_\sigma (\sigma^\dagger \sigma)^2 + \kappa_S (S^\dagger S)^2 + (\kappa_{H\sigma} \mathcal{H}^\dagger \mathcal{H} + \kappa_{\eta\sigma} \eta^\dagger \eta)(\sigma^\dagger \sigma) + (\kappa_{HS} \mathcal{H}^\dagger \mathcal{H} + \kappa_{\eta S} \eta)(S^\dagger S) \\
& + \kappa_{\sigma S} (\sigma^\dagger \sigma)(S^\dagger S) + m_H^2 \mathcal{H}^\dagger \mathcal{H} + m_\eta^2 \eta^\dagger \eta + m_\sigma^2 \sigma^\dagger \sigma + m_S^2 S^\dagger S + V_b.
\end{aligned} \tag{2}$$

We have introduced the extension of SM with global $U(1) \times Z_4$ symmetry and several additional fields. We introduced vector-like down-type quarks (D_L, D_R) , a pair of vector-like charged leptons (E_L, E_R) , and three right-handed singlet fermions $N_j, j = 1, 2, 3$. We introduced the doublet η , which later in this model can be chosen as the dark matter (DM) candidate for the lightest neutral component with Z_2 odd parity, and singlet-scalars σ and S which later be used as the linear combination of inflaton in this model. The SM particles do not have a charge in this global symmetry. This global $U(1)$ has a color anomaly similar to the KSVZ model for the strong CP problem [1, 2], then it can play the role of PQ symmetry. Please note that $\mathcal{H}^\dagger = \frac{1}{\sqrt{2}}(0 \ h^*)$ is the SM Higgs doublet. Also, we have d_{R_j} and e_{R_j} as the SM down-type quarks and charged leptons. We assume, all parameters in Lagrangian are real-positive and the dominant parts are up to dimension five with the cut-off scale of M_* . Lastly, we identify V_b as another potential, its terms are invariant under global symmetry but violate the S number.

$$V_b = \alpha(S^4 + S^{\dagger 4}) + \beta \sigma^\dagger \sigma (S^2 + S^{\dagger 2}) = \frac{1}{2} \tilde{S}^2 (\alpha \tilde{S}^2 \cos 4\rho + \beta \tilde{\sigma}^2 \cos 2\rho), \tag{3}$$

where we define $\sigma = \frac{\tilde{\sigma}}{\sqrt{2}} e^{i\theta}$ and $S = \frac{\tilde{S}}{\sqrt{2}} e^{i\rho}$. Please see fig. I for the corresponding charge under the symmetry we introduced.

	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)$	Z_4		$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)$	Z_4
D_L	3	1	$-\frac{1}{3}$	0	2	D_R	3	1	$-\frac{1}{3}$	2	0
E_L	1	1	-1	0	2	E_R	1	1	-1	2	0
σ	1	1	0	-2	2	S	1	1	0	0	2
N_k	1	1	0	1	1	η	1	2	$-\frac{1}{2}$	-1	1

TABLE I. The additional fields in respect to this model which charge under $[SU(3)_C \times SU(2)_L \times U(1)_Y] \times U(1) \times Z_4$

II. THE CP VIOLATION PHASES IN CKM AND PMNS MATRICES

To determine the Yukawa couplings of down-type quarks and charged leptons in the Lagrangian (1), the mass terms can be written by

$$\begin{pmatrix} \bar{f}_{L_i} & \bar{F}_L \end{pmatrix} \mathcal{M}_f \begin{pmatrix} f_{R_j} \\ F_R \end{pmatrix} + \text{h.c.}, \quad \mathcal{M}_f = \begin{pmatrix} m_{f_{ij}} & 0 \\ \mathcal{F}_{fj} & \mu_F \end{pmatrix} = \begin{pmatrix} h_{f_{ij}} h & 0 \\ (y_{fj} u e^{i\rho} + \tilde{y}_{fj} u e^{-i\rho}) & y_F w \end{pmatrix}. \quad (4)$$

The f and F correspond to $f = d, e$ and $F = D, E$ for down-type quarks and charge leptons, also \mathcal{M}_f is 4×4 mass matrix written in eq. (4). This mass matrix has a similar form [3]. Here, we suggest that global $U(1)$ symmetry works as the PQ symmetry and all parameters are assumed to be real, the $\arg(\det \mathcal{M}_f) = 0$ is fulfilled, whether the radiative effects are taken into account after the braking of PQ symmetry (see [4, 5]).

In this thesis, we assume that $\mu_F^2 + \mathcal{F}_f \mathcal{F}_f^\dagger \gg \mathcal{F}_f m_f^\dagger$. It can be obviously satisfied as we assume $u, w \gg h$. We can also have

$$B_f \simeq -\frac{A_f m_f \mathcal{F}_f^\dagger}{\mu_F^2 + \mathcal{F}_f \mathcal{F}_f^\dagger}, \quad C_f \simeq \frac{\mathcal{F}_f m_f^\dagger}{\mu_F^2 + \mathcal{F}_f \mathcal{F}_f^\dagger}, \quad D_f \simeq 1. \quad (5)$$

These approximately guarantee the unitarity of the matrix A_f . It is easy to find

$$A_f^{-1} \tilde{m}_f^2 A_f \simeq m_f m_f^\dagger - \frac{1}{\mu_F + \mathcal{F}_f \mathcal{F}_f^\dagger} (m_f \mathcal{F}_f^\dagger) (\mathcal{F}_f m_f^\dagger). \quad (6)$$

III. EFFECTIVE MODEL AT LOWER ENERGY REGION

In this section, we turn back to the Lagrangian eq. (I) and (2) and integrating out the heavy fields σ and S , as in this part we discuss the effective model in the lower energy region.

It is leaving us with the SM with extended lepton sector as it sounds in Scotogenic model referred in [6] which invariant under Z_2 symmetry. The Lagrangian is

$$\begin{aligned}
-\mathcal{L}_{scotogenic} = & \sum_{j=1}^3 \left[\frac{M_{N_j}}{2} \bar{N}_j^c N_j + \sum_{\alpha=1}^3 \tilde{h}_{\alpha j}^* \eta \bar{l}_{\alpha} N_j + \text{h.c.} \right] + \tilde{m}_H^2 \mathcal{H}^\dagger \mathcal{H} + \tilde{m}_\eta^2 \eta^\dagger \eta \\
& + \tilde{\lambda}_1 (\mathcal{H}^\dagger \mathcal{H})^2 + \tilde{\lambda}_2 (\eta^\dagger \eta)^2 + \tilde{\lambda}_3 (\mathcal{H}^\dagger \mathcal{H})(\eta^\dagger \eta) + \lambda_4 (\mathcal{H}^\dagger \eta)(\eta^\dagger \mathcal{H}) + \frac{\tilde{\lambda}_5}{2} [\sigma(\eta^\dagger \mathcal{H})^2 + \text{h.c.}] ,
\end{aligned} \tag{7}$$

which the redefined coupling constant (with tilde) are appearing after the symmetry breaking of $\tilde{\sigma}$ and \tilde{S} :

$$\begin{aligned}
\tilde{\lambda}_1 &= \lambda_1 - \frac{\kappa_{H\sigma}^2}{4\tilde{\kappa}_\sigma} - \frac{\kappa_{HS}^2}{4\tilde{\kappa}_S} + \frac{\kappa_{\sigma S} \kappa_{H\sigma} \kappa_{HS}}{4\tilde{\kappa}_\sigma \tilde{\kappa}_S}, & \tilde{\lambda}_2 &= \lambda_2 - \frac{\kappa_{\eta\sigma}^2}{4\tilde{\kappa}_\sigma} - \frac{\kappa_{\eta S}^2}{4\tilde{\kappa}_S} + \frac{\kappa_{\sigma S} \kappa_{\eta\sigma} \kappa_{\eta S}}{4\tilde{\kappa}_\sigma \tilde{\kappa}_S} \\
\tilde{\lambda}_3 &= \lambda_3 - \frac{\kappa_{H\sigma} \kappa_{\eta\sigma}}{2\tilde{\kappa}_\sigma} - \frac{\kappa_{HS} \kappa_{\eta S}}{2\tilde{\kappa}_S} + \frac{\kappa_{\sigma S} \kappa_{H\sigma} \kappa_{\eta S} + \kappa_{\sigma S} \kappa_{\eta\sigma} \kappa_{HS}}{4\tilde{\kappa}_\sigma \tilde{\kappa}_S}, & \tilde{\lambda}_5 &= \lambda_5 \frac{w}{M_*}.
\end{aligned} \tag{8}$$

In addition, the mass parameters in this low scale can be written as

$$\begin{aligned}
M_{N_j} &= y_{N_j} w \\
\tilde{m}_H^2 &= m_H^2 + \left(\kappa_{H\sigma} + \frac{\kappa_{HS} \kappa_{\sigma S}}{2\tilde{\kappa}_S} \right) w^2 + \left(\kappa_{HS} + \frac{\kappa_{HS} \kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} \right) u^2, \\
\tilde{m}_\eta^2 &= m_\eta^2 + \left(\kappa_{\eta\sigma} + \frac{\kappa_{\eta S} \kappa_{\sigma S}}{2\tilde{\kappa}_S} \right) w^2 + \left(\kappa_{\eta S} + \frac{\kappa_{\eta S} \kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} \right) u^2,
\end{aligned} \tag{9}$$

where in this model η has no VEV if $\tilde{m}_\eta^2 > 0$. We assume both \tilde{m}_H and \tilde{m}_η is in order of 0(1)TeV, parameters tuning are required for this.

As in the original scotogenic model, the neutrino mass is forbidden at tree level due to Z_2 symmetry but it can be generated through a 1-loop diagram. The result is given by

$$\mathcal{M}_{\alpha\beta}^\nu \simeq \sum_{j=1}^3 \tilde{h}_{\alpha j} \tilde{h}_{\beta j} \tilde{\lambda}_5 \Lambda_j \quad \text{with} \quad \Lambda = \frac{h^2}{8\pi} \frac{1}{M_{N_j}} \ln \frac{M_{N_j}^2}{M_\eta^2}, \tag{10}$$

where we supposed $M_\eta^2 = \tilde{m}_\eta^2 + (\tilde{\lambda}_3 + \lambda_4)h^2$ and $M_{N_j} \gg M_\eta$. As an example, we can adjust some parameters of neutrino Yukawa coupling as [7].

$$\tilde{h}_{ei} = 0, \quad \tilde{h}_{\mu i} = h_{\tau i} = h_i \quad (i = 1, 2); \quad \tilde{h}_{e3} = \tilde{h}_{\mu 3} = -\tilde{h}_{\tau 3} = h_3 \tag{11}$$

In this section, we will discuss the Inflation due to singlet scalars. It can be traced back to the similar problems mentioned in ref. [8–13]. For the remaining of this inflation part, we borrow the relevant method from ref. [14] and [15]. Straightforwardly, in this model, the action relevant to inflation is given by

$$S_J = \int d^4x \sqrt{-g} \left[-\frac{1}{2} M_p^2 R - \xi_\sigma \sigma^\dagger \sigma R - \tilde{\xi}_S S^\dagger S R + \partial^\mu \sigma^\dagger \partial_\mu \sigma + \partial^\mu S^\dagger \partial_\mu S - V(\sigma, S) \right]. \quad (12)$$

We impose the conformal transformation

$$g_{\mu\nu} = \Omega^2 g_{\mu\nu} \quad \Omega^2 = 1 + \frac{\xi_\sigma \tilde{\sigma}^2 + \tilde{\xi}_S \tilde{S}^2}{M_p^2}. \quad (13)$$

. Using the transformation, we can obtain the Action in the Einstein frame as

$$S_E = \int d^4x \sqrt{-g} \left[-\frac{1}{2} M_p^2 R_E + \frac{1}{2} \partial^\mu \phi_\sigma \partial_\mu \phi_\sigma + \frac{1}{2} \partial^\mu \phi_S \partial_\mu \phi_S + \frac{6\xi_\sigma \tilde{\xi}_S \tilde{\sigma} \tilde{S}}{\left[\left(\Omega^2 + \frac{6\xi_\sigma^2}{M_p^2} \tilde{\sigma} \right) \left(\Omega^2 + \frac{6\tilde{\xi}_S^2}{M_p^2} \tilde{S} \right) \right]^{1/2}} \partial^\mu \phi_\sigma \partial_\mu \phi_S - \frac{1}{\Omega^4} V(\tilde{\sigma}, \tilde{S}) \right], \quad (14)$$

where the superscript E represents the Einstein frame. In addition, three regions are containing the inflaton potential

$$V(\phi) = \begin{cases} \frac{\hat{\kappa}_S}{4\xi_S^2} M_p^4 \left[1 - \exp\left(-\sqrt{\frac{2}{3}} \frac{\phi}{M_p}\right) \right]^2 & \text{if } \phi > M_p \\ \frac{\hat{\kappa}_S}{6\xi_S^2} M_p^2 \phi^2 & \text{if } \frac{M_p}{\xi_S} < \phi < M_p \\ \frac{1}{4} \hat{\kappa} \phi^4, & \text{if } \phi < \frac{M_p}{\xi_S}, \end{cases} \quad (15)$$

The first region ($\phi > M_p$) correspond to the inflationary phase, the inflation ends when $\phi \simeq M_p$. After the end of inflation, if $\tilde{\xi}_S > 1$, the inflation will reach the quadratic potential depicted in the second region ($M_p/\tilde{\xi}_S < \phi < M_p$).

It is necessary to calculate the slow roll parameters as

$$\epsilon \equiv \frac{1}{2} M_p \left(\frac{V'}{V} \right)^2 = \frac{8M_p^4}{b\tilde{\xi}_S (1 + 6\tilde{\xi}_S/b) \tilde{S}^4}, \quad \eta \equiv M_p^2 \frac{V''}{V} = -\frac{8M_p^2}{b (1 + 6\tilde{\xi}_S/b) \tilde{S}^2}. \quad (16)$$

The number of e-folds \mathcal{N}_k with scale k exits the Horizon to the end of inflation can be calculated as

$$\mathcal{N}_k = \frac{1}{M_p} \int_{\phi_{end}}^{\phi_k} \frac{V}{V'} d\phi = \frac{1}{8M_p^2} (b + 6\tilde{\xi}_S) (\tilde{S}_k^2 - \tilde{S}_{end}^2) - \frac{3}{4} \ln \frac{M_p^2 + \tilde{\xi}_S \tilde{S}_k^2}{M_p^2 + \tilde{\xi}_S \tilde{S}_{end}^2}, \quad (17)$$

With this, we have another approximate relation, there are: $\epsilon \simeq \frac{3}{4\mathcal{N}_K}$ and $\eta \simeq -\frac{1}{\mathcal{N}_K}$. The potential during the end of inflation is approximated to be $V(\phi) \simeq 0.072 \frac{\hat{\kappa}_S}{\xi_S^2} M_p^4$.

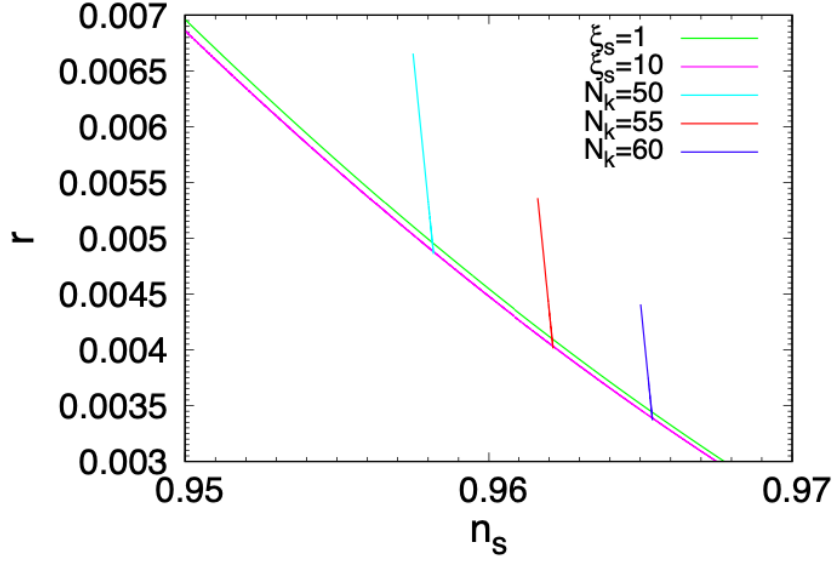


FIG. 1. We varied the coupling constant $\hat{\kappa} \sim 10^{-7} - 10^{-10}$ with the values of n_s and r can be read-off by the intersection of fixed $\tilde{\xi}_S$ or \mathcal{N}_k .

The scalar power spectrum can be written as

$$\mathcal{P}(k) = A_s \left(\frac{k}{k_*} \right)^{n_s-1} \quad A_s = \frac{V^3}{12\pi^2 M_p^6 V'^2} = \frac{V}{24\pi^2 M_p^4 \epsilon} \Big|_{k_*}. \quad (18)$$

If we used the Planck data $A_s = (2.101_{-0.034}^{+0.031}) \times 10^{-9}$ at $k^* = 0.05 \text{Mpc}^{-1}$ [16]. We find the constraint

$$\hat{\kappa}_S \simeq 4.13 \times 10^{-10} \tilde{\xi}_S^2 \left(\frac{60}{\mathcal{N}_{k^*}} \right)^2. \quad (19)$$

As we already have all requirements to calculate the spectral index and the tensor-to-scalar ratio respectively

$$r = 16\epsilon \quad n_s = 1 - 6\epsilon + 2\eta. \quad (20)$$

Since $\hat{\kappa}_S$ is free parameter in our model, using the the fixed values of $\tilde{\xi}_S$ and \mathcal{N}_\parallel , we choose the range $10^{-10} \leq \hat{\kappa}_S \leq 10^{-7}$. The constraint of CMB (19) can be obtained in the intersection points of the fixed $\tilde{\xi}_S$ and \mathcal{N}_k (see fig. 1).

IV. END OF INFLATION: PREHEATING AND REHEATING

During the end of Inflation, the friction ($3H\dot{\phi}$) is getting smaller, and comparable with other terms. With this, we can write the Klein-Gordon equation as

$$\ddot{\phi} + \frac{dV(\phi)}{d\phi} \simeq 0. \quad (21)$$

As we already stated, the quadratic potential is neglected due to the reason we mentioned before. Thus the quartic potential plays a substantial role to drain the inflaton energy. If we introduce the dimensionless conformal time τ as $a\tau = \sqrt{\hat{\kappa}_S}\phi_{end}$ and also we redefine the field $f = \frac{a\phi}{\phi_{end}}$, eq. (21). Hence, it can be approximated by

$$\frac{d^2 f}{d\tau^2} + f^3 = 0. \quad (22)$$

The solution of this equation is belong to Jacobi Elliptic function $f(\tau) = \text{cn}\left(\tau - \tau_i, \frac{1}{\sqrt{2}}\right)$ [14, 15]. The last equation is obtained by using redefinition of conformal time as

$$a(\tau) = \frac{\phi_{end}}{2\sqrt{3}M_p}\tau \quad \tau = 2(3\hat{\kappa}_S M_p^2)^{1/4}\sqrt{t}. \quad (23)$$

In the purpose of our discussion, we will recall some terms in Lagrangian (I) and (2) which correspond to the the field $\tilde{\sigma}$ and \tilde{S} and substitute both fields with ϕ , we obtain

$$\begin{aligned} & \left[-\frac{y_D}{\sqrt{2}} \frac{\kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} \phi \bar{D}_L \bar{D}_R - \frac{y_E}{\sqrt{2}} \frac{\kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} \phi \bar{E}_L \bar{E}_R - \sum_{j=1}^3 \left\{ \frac{1}{\sqrt{2}} (y_{d_j} e^{i\rho} + \tilde{y}_{d_j} e^{-i\rho}) \phi \bar{D}_L d_{R_j} \right. \right. \\ & \quad \left. \left. + \frac{1}{\sqrt{2}} (y_{e_j} e^{i\rho} + \tilde{y}_{e_j} e^{-i\rho}) \phi \bar{E}_L e_{R_j} + \frac{y_{N_j}}{2\sqrt{2}} \frac{\kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} \phi \bar{N}_j^c N_j \right\} + \text{h.c} \right] \\ & + \frac{1}{2} (\kappa_{\sigma S} \mathcal{H}^\dagger \mathcal{H} + \kappa_{\eta S} \eta^\dagger \eta) \phi^2 - \frac{1}{2} \frac{\kappa_{\sigma S}}{2\tilde{\kappa}_\sigma} (\kappa_{H\sigma} \mathcal{H}^\dagger \mathcal{H} + \kappa_{H\sigma} \eta^\dagger \eta) \phi^2. \end{aligned} \quad (24)$$

Thus, the mass term from eq. (24) can be extracted explicitly

$$\begin{aligned} M_{N_j} &= \frac{y_{N_j}}{\sqrt{2}} \frac{|\kappa_{\sigma S}|}{2\tilde{\kappa}_\sigma} \phi & \tilde{M}_F &= \frac{\phi}{\sqrt{2}} \left[\sum_{j=1}^3 (y_{f_j}^2 + \tilde{y}_{f_j}^2) + y_F^2 \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_\sigma^2} \right]^{1/2} \\ m_H^2 &= \frac{1}{2} \left(\kappa_{HS} + \frac{|\kappa_{\sigma S}|}{2\tilde{\kappa}_\sigma} \kappa_{H\sigma} \right) \phi^2 & m_\eta^2 &= \frac{1}{2} \left(\kappa_{\eta S} + \frac{|\kappa_{\sigma S}|}{2\tilde{\kappa}_\sigma} \kappa_{\eta\sigma} \right) \phi^2. \end{aligned} \quad (25)$$

Where $F = D$ or E should be regarded similar with $f = d$ or e . In addition, we will recall back and and split \tilde{S}^2 to be $\tilde{S}^2 = S_\parallel^2 + S_\perp^2$, also the same way with $\tilde{\sigma}$, we finally have

$$\frac{\tilde{\kappa}_S}{4} (S_\parallel^2 + S_\perp^2 - u^2)^2 + \frac{\kappa_{\sigma S}}{4} (S_\parallel^2 + S_\perp^2 - u^2) (\sigma_\parallel^2 + \sigma_\perp^2 - w^2) + \frac{\tilde{\kappa}_\sigma}{4} (\sigma_\parallel^2 + \sigma_\perp^2 - w^2)^2. \quad (26)$$

Using the last potential, we can obtain the masses of each component:

$$\begin{aligned}
m_{S_{\parallel}}^2 &= \frac{\partial^2 V}{\partial S_{\parallel}^2} = \left(3\tilde{\kappa}_S S_{\parallel}^2 + \frac{\kappa_{\sigma S}}{2}(\sigma_{\parallel}^2 + \sigma_{\perp}^2 - w^2) \right) = \left(3\tilde{\kappa}_S S_{\parallel}^2 - \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}}(S_{\parallel}^2 + S_{\perp}^2 - u^2) \right) \\
&\simeq \left(3\tilde{\kappa}_S S_{\parallel}^2 - \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}} S_{\parallel}^2 \right) = \left(3\tilde{\kappa}_S S_{\parallel}^2 - \frac{\kappa_{\sigma S}^2}{2\tilde{\kappa}_{\sigma}} S_{\parallel}^2 + \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}} S_{\parallel}^2 \right) \\
&= \left(3\hat{\kappa}_S S_{\parallel}^2 - \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}} S_{\parallel}^2 \right) = \left(3\hat{\kappa}_S^2 + \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}} \right) \phi^2.
\end{aligned} \tag{27}$$

where we have used $\tilde{\sigma} = \sqrt{\frac{|\kappa_{\sigma S}|}{2\tilde{\kappa}_S}} \tilde{S}$ and $S_{\parallel} \gg S_{\perp}, u$, since in this part we regard S_{\parallel} as inflaton. Using the same method, as we skipped their calculation for better reason, the other masses can be obtained:

$$m_{S_{\perp}}^2 = \hat{\kappa}_S \phi^2, \quad m_{\sigma_{\parallel}}^2 = \left(|\kappa_{\sigma S}| + \frac{\kappa_{\sigma S}^2}{4\tilde{\kappa}_{\sigma}} \right) \phi^2, \quad \text{and} \quad m_{\sigma_{\perp}}^2 = \frac{\kappa_{\sigma S}}{4\tilde{\kappa}_{\sigma}} \phi^2. \tag{28}$$

It is clear, the masses of particles are ϕ dependent, so the decay channel is opened once $\phi \simeq 0$. Hence, the particles are produced during zero crossings. Here we write the equation of oscillation for the self-production of the inflaton

$$\frac{d^2 \Phi_k}{d\tau^2} + \omega_k^2 \Phi_k = 0, \quad \omega^2 = \bar{k}^2 + 3f(\tau)^2, \tag{29}$$

and the equation of the created particles

$$\frac{d^2 F_k}{d\tau^2} + \tilde{\omega}_k^2 F_k = 0, \quad \tilde{\omega}^2 = \bar{k}^2 + \frac{g_{\psi}}{\hat{\kappa}_S} f(\tau)^2, \tag{30}$$

where we have used the rescaled variables

$$\Phi_k = \frac{a\phi_k}{\phi_{end}}, \quad F_k = \frac{a\psi_k}{\phi_{end}}, \quad \bar{k} = \frac{ak}{\phi_{end}\sqrt{\hat{\kappa}_S}}. \tag{31}$$

Here $f(\tau)$ is the solution of Eq. (22). Both Φ_k and F_k shown the exponential behavior $\propto e^{\mu_k \tau}$, where μ_k represents the characteristic exponent [14, 17]. This μ_k is determined by the ratio of $g_{\psi}/\hat{\kappa}_S$. The number density of produced particle of each species ψ can be written as

$$n_k^{\psi} = \frac{\tilde{\omega}_k}{2\hat{\kappa}_S} \left(\frac{|F_k|^2}{\tilde{\omega}_k} + |F_k|^2 \right) - \frac{1}{2}. \tag{32}$$

We then parametrize the coupling g_{ψ} into 5 groups

$$(A) \frac{g_{\sigma_{\parallel}}}{\hat{\kappa}_S} \gg 1, \quad (B) \frac{g_{S_{\parallel}}}{\hat{\kappa}_S} = 3, \quad (C) \frac{g_{S_{\perp}}}{\hat{\kappa}_S} = 1, \quad (D) \frac{g_{\sigma_{\perp}}}{\hat{\kappa}_S} \ll 1, \quad (E) \frac{g_H}{\hat{\kappa}_S}, \frac{g_{\eta}}{\hat{\kappa}_S} > 1. \tag{33}$$

We conclude here, the (A)-(D) cases are well constrained by the composition of the inflationary model, meanwhile, (E) are not.

The momentum distribution of produced particle ψ -species through the one-zero crossing of inflaton ϕ is

$$n_{\bar{k}}^{\psi} = e^{2\mu_k \frac{\tau_0}{2}} = e^{-\left(\frac{\bar{k}}{\bar{k}_c}\right)^2}, \quad \bar{k}_c^2 = \sqrt{\frac{g_{\psi}}{2\pi^2, \hat{\kappa}_S}} \quad (34)$$

where τ_0 is the inflaton period. The resonance is efficient for $\bar{k} < \bar{k}_c$, then the particle number density of species ψ can be calculated via

$$n^{\psi} = \int \frac{d^3\bar{k}}{(2\pi)^3} n_{\bar{k}}^{\psi} = \int \frac{d^3\bar{k}}{(2\pi)^3} e^{-\left(\frac{\bar{k}}{\bar{k}_c}\right)^2} = \frac{\bar{k}_c^3}{8\pi^{3/2}}. \quad (35)$$

The energy transfer from inflaton to relativistic particles happens from the decay of \mathcal{H} and η . Thus, the relativistic particles are created by indirect ones. The $\mathcal{H} \rightarrow \bar{q}t$ decay process with top Yukawa coupling h_t while $\eta \rightarrow \bar{l}N$ with neutrino Yukawa coupling h_j . During the oscillation period, the induced mass η can be larger than N_j . The decay width of $\psi = \mathcal{H}, \eta$ can be written by

$$\bar{\Gamma}_{\psi} = \frac{c_{\psi} y_{\psi}^2}{8\pi} \bar{m}_{\psi} \quad \bar{m}_{\psi} = \frac{am_{\psi}}{\phi_{end} \sqrt{\hat{\kappa}_S}} = \sqrt{\frac{g_{\psi}}{\hat{\kappa}_S}} f(\tau), \quad (36)$$

where c_{ψ} is internal degrees of freedom, $c_H = 3$ and $c_{\eta} = 1$. The Yukawa coupling y_{ψ} represents $y_H = h_t$ and $y_{\eta} = h_j$. For $\bar{\Gamma}_{\psi}^{-1} < \tau_0/2$ is satisfied with $g_{\psi} > 4 \times 10^{-7} \left(\frac{\hat{\kappa}_S}{10^{-8}}\right)$, the produced ψ -species decays to the relativistic fermions are finished before next zero-crossing [18]. If we fix $\tau = 0$ at the first zero-crossing, we can approximate $f(\tau) = \sin(cf_0\tau)$. The energy transferred by ψ decay can be written by

$$\delta\bar{\rho}_r = \int_0^{\tau_0/2} d\tau \bar{\Gamma}_{\psi} \bar{m}_{\psi} \bar{n}_{\psi} e^{-\int_0^{\tau} \bar{\Gamma}_{\psi} \tau'} = \frac{1}{8\pi^{3/2} (2\pi^2)^{3/4}} \left(\frac{g_{\psi}}{\hat{\kappa}_S}\right)^{5/4} Y(f_0, \gamma_{\psi}). \quad (37)$$

where γ_{ψ} and $Y(f_0, \gamma_{\psi})$ are defined as

$$\gamma_{\psi} = \frac{c_{\psi} y_{\psi}^2}{8\pi c} \sqrt{\frac{g_{\psi}}{\hat{\kappa}_S}}, \quad Y(f_0, \gamma_{\psi}) = c\gamma_{\psi} \int_0^{\tau_0/2} d\tau f_0^2 \sin^2(cf_0\tau) e^{-2\gamma_{\psi} \sin^2(\frac{cf_0\tau}{2})}, \quad (38)$$

here we used $c = 2\pi/\tau_0$. The energy density which is transferred to the light particles is accumulated for each zero-crossing, it can be approximated using average value of τ as

$$\bar{\rho}_r = \frac{2\tau}{\tau_0} \delta\bar{\rho}_r = 6.5 \times 10^{-4} \left(\frac{g_{\psi}}{\hat{\kappa}_S}\right)^{5/4} Y(f_0, \gamma_{\psi}) \tau, \quad (39)$$

where we assumed f_0 to be constant. The total energy density of the inflation energy $\bar{\rho}_{\phi}$ and transfer energy $\bar{\rho}_r$ are conserved, the reheating temperature can be found by using relation $\bar{\rho}_{\phi} = \bar{\rho}_r$. We found

$$\frac{1}{4\hat{\kappa}_S} \left(\frac{\sqrt{\hat{\kappa}_S} \phi_{end}}{a}\right)^4 = \frac{\pi^2}{30} g_* T_R^4, \quad (40)$$

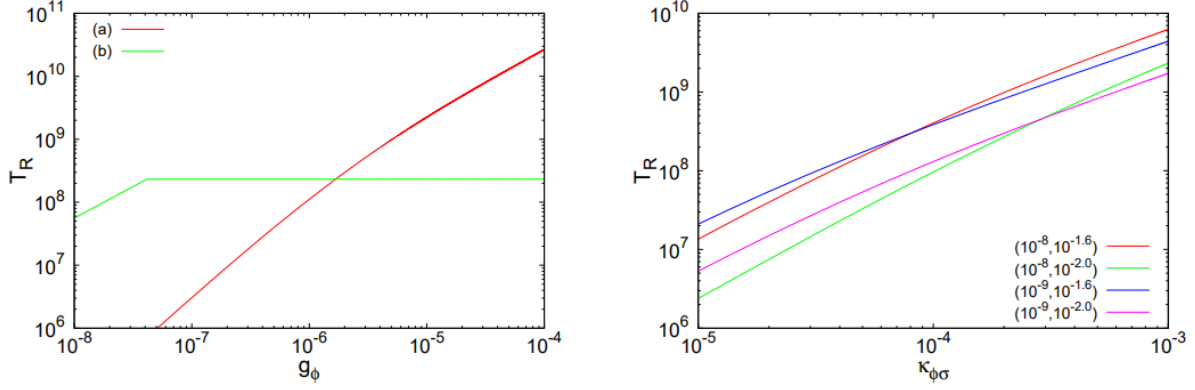


FIG. 2. **left:** the comparison of the reheating temperature T_R in (a) preheating and (b) perturbative process, where the assumed perturbative process is $h \rightarrow \bar{q}t$. **right:** the reheating temperature due to σ . We fix $\tilde{\kappa}_\sigma = (10^{4.5}, 10^{5.3})$ and vary the parameters $\tilde{\kappa}_S, \kappa_{\sigma S}/\tilde{\kappa}_\sigma$.

we used $\bar{\rho}_\phi = \frac{1}{4\tilde{\kappa}_S}$ and $g_* = 130$. Using relation (23) and (39) to (40) we obtain

$$T_R = 5.9 \times 10^{15} g_\psi^{5/4} Y(f_0, \gamma_\psi) \text{GeV}. \quad (41)$$

V. LEPTOGENESIS

In the ordinary seesaw model, the neutrino mass is generated by Yukawa interaction $h_{\alpha j} \bar{l}_\alpha \eta N_j$. The production of N_1 , the lightest right-handed-neutrino, in the thermal bath, is due to the scattering process of $\tilde{D}_L D_R, \tilde{E}_L E_R \rightarrow N_1 N_1$ mediated by $\tilde{\sigma}$ if both fermions are in thermal equilibrium. The conditions $T > \tilde{M}_F, M_{N_1}$ and $\Gamma_{FF} \simeq H$ must be satisfied, where Γ_{FF} is the reaction rate of the scattering. The estimation of the temperature using relation of $\Gamma_{FF} \simeq H(T)$ gives

$$T \simeq 5.8 \times 10^8 \left(\frac{y_F}{10^{-1.2}} \right)^2 \left(\frac{y_{N_1}}{10^{-2}} \right)^2 \text{GeV}. \quad (42)$$

Thus, $T > \tilde{M}_F, M_{N_1}$ if y_F and y_{N_1} satisfied.

after N_1 is created by scattering of extra fermions, the decay product $l_\alpha \eta^\dagger$ is expected by suppressed Yukawa coupling $h_{\alpha 1}$. After the washout process which is labeled as frozen out, the decay occurs and so the lepton number asymmetry which is generated can be efficiently converted to baryon number asymmetry via sphaleron processes. We can investigate this case by solving Boltzmann equation for Y_{N_1} and $Y_L (\equiv Y_l - Y_{\bar{l}})$. $Y_\psi = \frac{n_\psi}{s}$, with number

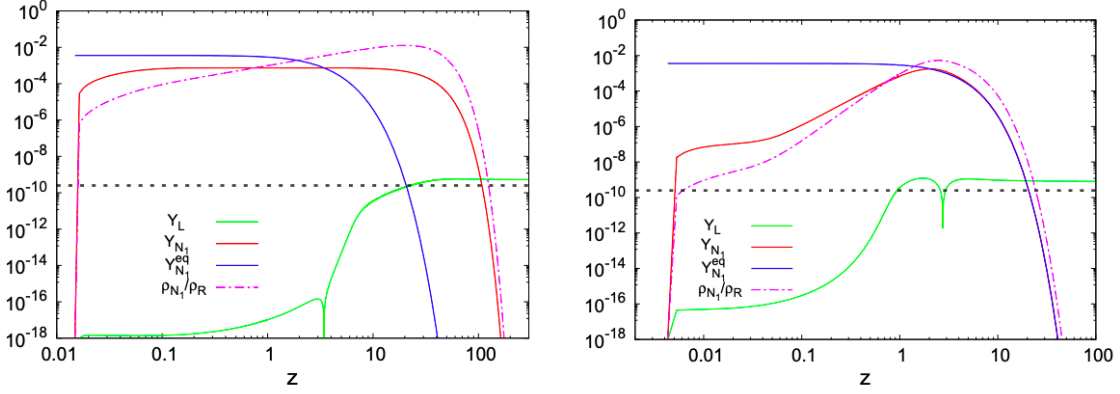


FIG. 3. Both Figures correspond to the evolution of Y_L and Y_{N_1} . Case (I) is depicted in left-panel with $\kappa_{H\sigma} = 10^{-4}$ and $\kappa_{HS} = \kappa_{\eta S} = 0$ Case (II) is depicted in right-panel for $\kappa_{HS} = \kappa_{\eta S} = 10^{-6}$ while $\kappa_{H\sigma} = \kappa_{\eta\sigma} = 0$. Please note, the other parameters depicted in the the text. The initial condition is simply depicted as $Y_L = Y_{N_1} = 0$ at $z = z_R$ while ρ_{N_1}/ρ_R depicted as ratio of energy density of N_1 and radiation.

density n_ψ and entropy density s . Solving the Boltzman equation, we see the plot in fig. 3.

VI. DARK MATTER AND ISOCURVATURE FLUCTUATIONS

In this model, we propose 2 dark matters candidate. The first is the lightest component of η with Z_2 odd parity [19–23]. If we assume η_R is the main component of the dark matter, the dark matter abundance and its direct search can be preserved if $\tilde{\lambda}_3$ and $|\lambda_4|$ take the suitable values. However, these parameters may affect the perturbativity of the quartic coupling due to radiative corrections, but we can safely stay away from this problem in certain regions. The other candidate for dark matter is Axion. This happens if $f_a \sim 10^{11}\text{GeV}$. This case is can be inferred to the case (II). The PQ symmetry is spontaneously broken during inflation. The Axion is depicted as the phase θ in $\sigma = \frac{1}{\sqrt{2}}\tilde{\sigma}e^{i\theta}$ and its potential is flat during inflation. The Axion gets the quantum fluctuation $\delta A = \left(\frac{H}{2\pi}\right)^2$ and causes the isocurvature fluctuation and affects the CMB amplitude [24–26].

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- [1] J. E. Kim, Weak-interaction singlet and strong CP invariance, *Phys. Rev. Lett.* **43**, 103 (1979).
 - [2] M. A. Shifman, A. Vainshtein, and V. I. Zakharov, Can confinement ensure natural cp invariance of strong interactions?, *Nuclear Physics B* **166**, 493 (1980).
 - [3] L. Bento, G. C. Branco, and P. A. Parada, A minimal model with natural suppression of strong cp violation, *Physics Letters B* **267**, 95 (1991).
 - [4] D. Suematsu, Pati-salam unification with a spontaneous c p violation, *Physical Review D* **100**, 055019 (2019).
 - [5] D. Suematsu, Low scale leptogenesis in a model with promising cp structure, *The European Physical Journal C* **81**, 1 (2021).
 - [6] E. Ma, Verifiable radiative seesaw mechanism of neutrino mass and dark matter, *Physical Review D* **73**, 10.1103/physrevd.73.077301 (2006).
 - [7] D. Suematsu, T. Toma, and T. Yoshida, Reconciliation of cdm abundance and $\mu \rightarrow e \gamma$ in a radiative seesaw model, *Physical Review D* **79**, 093004 (2009).
 - [8] R. N. Lerner and J. McDonald, Gauge singlet scalar as inflaton and thermal relic dark matter, *Physical Review D* **80**, 123507 (2009).
 - [9] R. N. Lerner and J. McDonald, Distinguishing higgs inflation and its variants, *Physical Review D* **83**, 123522 (2011).
 - [10] R. H. Budhi, S. Kashiwase, and D. Suematsu, Inflation in a modified radiative seesaw model, *Physical Review D* **90**, 113013 (2014).
 - [11] R. H. Budhi, S. Kashiwase, and D. Suematsu, Constrained inflaton due to a complex scalar, *Journal of Cosmology and Astroparticle Physics* **2015** (09), 039.
 - [12] R. H. Budhi, S. Kashiwase, and D. Suematsu, Inflation due to a nonminimal coupling of singlet scalars in the radiative seesaw model, *Physical Review D* **93**, 013022 (2016).
 - [13] S. Kashiwase and D. Suematsu, Lepton number asymmetry via inflaton decay in a modified radiative seesaw model, *Physics Letters B* **749**, 603 (2015).
 - [14] P. B. Greene, L. Kofman, A. Linde, and A. A. Starobinsky, Structure of resonance in preheating after inflation, *Physical Review D* **56**, 6175 (1997).
 - [15] G. Ballesteros, J. Redondo, A. Ringwald, and C. Tamarit, Standard model—axion—seesaw—higgs portal inflation. five problems of particle physics and

- cosmology solved in one stroke, *Journal of Cosmology and Astroparticle Physics* **2017** (08), 001–001.
- [16] Y. Akrami, F. Arroja, M. Ashdown, J. Aumont, C. Baccigalupi, M. Ballardini, A. J. Banday, R. Barreiro, N. Bartolo, S. Basak, *et al.*, Planck 2018 results-x. constraints on inflation, *Astronomy & Astrophysics* **641**, A10 (2020), 1807.06209 [astro-ph.CO].
 - [17] L. Kofman, A. Linde, and A. A. Starobinsky, Towards the theory of reheating after inflation, *Physical Review D* **56**, 3258–3295 (1997).
 - [18] G. Felder, L. Kofman, and A. Linde, Instant preheating, *Physical Review D* **59**, 123523 (1999).
 - [19] S. Kashiwase and D. Suematsu, Baryon number asymmetry and dark matter in the neutrino mass model with an inert doublet, *Physical Review D* **86**, 053001 (2012).
 - [20] S. Kashiwase and D. Suematsu, Leptogenesis and dark matter detection in a tev scale neutrino mass model with inverted mass hierarchy, *The European Physical Journal C* **73**, 1 (2013).
 - [21] L. L. Honorez, E. Nezri, J. F. Oliver, and M. H. Tytgat, The inert doublet model: an archetype for dark matter, *Journal of Cosmology and Astroparticle Physics* **2007** (02), 028.
 - [22] T. Hambye, F.-S. Ling, L. L. Honorez, and J. Rocher, Scalar multiplet dark matter, *Journal of High Energy Physics* **2009**, 090 (2009).
 - [23] S. Andreas, M. H. Tytgat, and Q. Swillens, Neutrinos from inert doublet dark matter, *Journal of Cosmology and Astroparticle Physics* **2009** (04), 004.
 - [24] M. Beltran, J. Garcia-Bellido, and J. Lesgourgues, Isocurvature bounds on axions revisited, *Physical Review D* **75**, 103507 (2007).
 - [25] M. P. Hertzberg, M. Tegmark, and F. Wilczek, Axion cosmology and the energy scale of inflation, *Physical Review D* **78**, 083507 (2008).
 - [26] M. Fairbairn, R. Hogan, and D. J. Marsh, Unifying inflation and dark matter with the peccei-quinn field: observable axions and observable tensors, *Physical Review D* **91**, 023509 (2015).

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。）

Inflation and Reheating with The Singlet Scalars Related to CP Violation（CP の破れに關与する 1 重項スカラー場によるインフレーションと再加熱）2. 論文提出者 (1) 所 属 数物科学 専攻(2) 氏 名 ノーマ シディック リスディアント
Norma Sidik Risdianto

3. 審査結果の要旨（600～650 字）

初期宇宙において高温高密度のビッグバン宇宙に至る以前にインフレーションと呼ばれる指数関数的な膨張が存在したことが宇宙背景放射の観測により示唆されている。この膨張は通常インフラトンと呼ばれるスカラー場により引き起こされるとされ、インフラトンの解明は素粒子標準模型の拡張を考える際の重要な切り口となり得る。素粒子標準模型の構成要素であるヒッグス場は有力な候補と考えられているがいくつかの問題も指摘されており、標準模型の抱える問題と密接に関連した別のインフラトン候補の探索が望まれている。このような視点から申請者は、素粒子標準模型の重要な特徴の一つである CP 対称性の破れをもたらすクォークやレプトンの湯川結合定数の複素数位相の起源に着目する。

本博士論文では、スカラー場による自発的 CP の破れを、新たに導入されたフェルミオンとクォーク・レプトンとの混合を通して湯川結合の複素位相として発現させる拡張模型に注目し、自発的 CP の破れを引き起こすスカラー場をインフラトンとして採用することを提案し、その現象論的特徴を検討している。具体的には、湯川結合の複素位相の導出を可能にするという条件の下で (1)引き起こされるインフレーションの特徴、(2)インフレーション後に期待される再加熱現象、(3)ニュートリノ質量や暗黒物質残存量などに関する模型の予言について解析している。さらに、新たに導入されたフェルミオンの存在により、宇宙のバリオン数を十分に生成するために要求される右巻きニュートリノの質量下限を引き下げることが可能となることも指摘している。これらの内容は博士論文に値するものであると判定した。

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 (学 術)