Right-Handed Sneutrino Curvaton and non-Gaussianity

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Abstract

In this paper, we explore the parameter space for a Right-Handed (RH) sneutrino curvaton that can generate large non-Gaussianity without assuming any particular inflation sector. The mass of the RH sneutrino is suggested from a discussion on the initial condition of the curvaton field. It is shown that a small Yukawa coupling is generally required for a successful RH sneutrino curvaton. However, the Yukawa coupling can be larger if we consider the braneworld scenario. Some general discussion about the spectral index in curvaton scenario is also provided.

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1 Introduction

As we can see from observation, the Cosmic Microwave Background (CMB) is very isotropic but with small ($\sim O(10^{-5})$) temperature fluctuations. The advantage of single-field slow roll inflation models is that an inflaton field can produce inflation and also provide primordial density perturbation as the seeds of structure formation and CMB temperature fluctuation from the quantum fluctuation of the inflaton field during inflation. However, the model of inflation is highly restricted by the requirement of the right amount of primordial density (curvature) perturbations. There is one way to liberate the model of inflation if the job of creating curvature perturbation is done by another field which is called curvaton \cite{1,2,3,4}. If the curvaton field is light (smaller than the Hubble parameter), it can produce an almost scale invariant quantum fluctuation during inflation. However, by definition the energy density of our universe is dominated by the inflaton field during inflation, therefore the curvaton cannot produce curvature perturbation during inflation. The job must be done after inflation. If the curvaton decays after inflation, during the oscillation of the curvaton field (described as nonrelativistic matter), the universe will be dominated by radiation
after inflaton decay and the relative energy density of the curvaton is growing and can produce right amount of curvature perturbation. Here we can qualitatively know that the decay rate of a successful curvaton must be low.

We can distinguish between the cases that curvature perturbation is coming from inflaton or curvaton by investigating the non-Gaussianity of CMB (for a review of non-Gaussianity see Ref. [5]). The non-Gaussianity from single-field slow roll inflation is very small [6, 7, 8, 9] but it can be large from the curvaton scenario [10, 11, 12]. Conventionally, non-Gaussianity in curvaton scenario can be described by the non-linearity parameter $f_{NL}$, which takes the form

$$\zeta = \zeta_g + \frac{3}{5} f_{NL} \zeta_g^2 + \cdots,$$

where $\zeta$ is the curvature perturbation in the uniform density slice and $\zeta_g$ denotes the Gaussian part of $\zeta$. Currently the upper bound of $f_{NL}$ is roughly given by (2 − $\sigma$ range) [13, 14, 15]

$$f_{NL} \lesssim 100.$$

(2)

In the near future, the Planck satellite [16] will reduce the upper bound to $f_{NL} \lesssim 5$ if non-Gaussianity is not detected. Therefore, we will consider

$$10 \lesssim f_{NL} \lesssim 100,$$

(3)

which can be tested in the near future. We refer this range as large non-Gaussianity.

The possibility of using right-handed (RH) sneutrino as a curvaton was considered in [17, 18, 19, 20, 21]. However, the parameter space for generating large non-Gaussianity was not explored. A specific application of RH sneutrino to D-term hybrid inflation and non-Gaussianity was investigated in [22]. In this paper, we explore the parameter space in more generally settings without assuming any particular inflation model.

The paper is organized as follows. In Sec. 2 we present the formalism and describe non-Gaussianity generated in our RH sneutrino curvaton scenario. In Sec. 3 the initial condition of the curvaton field is discussed. This may suggest the Right-Handed sneutrino mass. In Sec. 4, we discuss the spectral index in the curvaton scenario. In Sec. 5 we consider the curvaton on a brane, and show that the Yukawa coupling can be larger in this case. Sec. 6 is our conclusion.

## 2 RH Sneutrino as a Curvaton

The superpotential of the mass eigenstate of the RH neutrino, $\Phi$, is given by

$$W_{\nu} = \lambda_{\nu} \Phi H_u L + \frac{m_\Phi^2}{2},$$

(4)

where $\Phi$ is the RH neutrino superfield, $H_u$ and $L$ are the MSSM Higgs and lepton doublet superfields, and $m$ is the RH neutrino mass. This gives the potential of right-handed sneutrino $\sigma$ as follows

$$V(\sigma) = \frac{1}{2} m^2 \sigma^2.$$  

(5)

The decay rate for RH sneutrino is

$$\Gamma = \frac{\lambda_{\nu}^2}{4\pi} m.$$  

(6)
In the following we will consider the case where the RH sneutrino is the curvaton. During inflation we require $m \ll H$ in order to have $\sigma$ slow-rolling, which means the field value can be taken as a constant during inflation.

The amplitude of quantum fluctuation of the curvaton field in a quasi-de Sitter space is given by

$$\delta \sigma = \frac{H^*}{2\pi},$$  \hspace{1cm} (7)

where $*$ denotes the epoch of horizon exit during inflation. The curvature perturbation generated from curvaton is given by

$$P_{\zeta_{\sigma}}^{1/2} = \frac{1}{3} \Omega_{\sigma,D} \frac{H^*}{\sigma_*}$$  \hspace{1cm} (8)

where

$$\Omega_{\sigma,D} \equiv \left( \frac{\rho_{\sigma}}{\rho_{\text{tot}}} \right)_D$$  \hspace{1cm} (9)

is the density fraction of the curvaton density $\rho_{\sigma}$ relative to the total density of the universe $\rho_{\text{tot}}$ at the time of curvaton decay, denoted by $D$, and $\zeta_{\sigma}$ is the curvature perturbation of the curvaton field $\sigma$. The amount of non-Gaussianity is characterized by the nonlinear parameter $f_{NL}$ given by

$$f_{NL} = \frac{5}{4\Omega_{\sigma,D}}.$$  \hspace{1cm} (10)

This equation is valid only when $\Omega_{\sigma,D} \ll 1$.

If we assume that at the time $t_o$ of curvaton oscillation with energy density $\rho_{\sigma}(t_o) = m^2 \sigma_*^2/2$, the universe is dominated by radiation (the decay products of inflaton) with energy density $\rho_R(t_o) = 3m^2 M_P^2$, where $M_P \equiv 2.4 \times 10^{18}$ GeV is the reduced Planck mass. Note that in order to have $\rho_{\sigma}(t_o) \ll \rho_R(t_o)$, we require $\sigma_* \ll M_P$. At the time of curvaton decay $t_D$, the energy density of the universe is given by $\rho_R(t_D) = 3\Gamma^2 M_P^2 = \rho_R(t_o)(a(t_o)/a(t_D))^4$. Therefore $a(t_D)/a(t_o) = (m/\Gamma)^{1/2}$ and $\Omega_{\sigma,D}$ is given by

$$\Omega_{\sigma,D} = \frac{\rho_{\sigma}(t_o) a(t_D)}{\rho_R(t_o) a(t_o)} = \frac{\sigma_*^2}{6M_P^2} \left( \frac{m}{\Gamma} \right)^{1/2} = \frac{\sigma_*^2}{6M_P^2} \frac{\sqrt{4\pi}}{\lambda}$$  \hspace{1cm} (11)

where $m$ is the Right-Handed sneutrino mass and $\Gamma$ its decay rate. If at the time of curvaton oscillation, the universe is dominated by oscillating inflaton field, we have

$$\Omega_{\sigma,D} = \frac{\sigma_*^2}{6M_P^2} \left( \frac{\Gamma_d}{\Gamma} \right)^{1/2} = \frac{\sigma_*^2}{6M_P^2} \frac{\sqrt{4\pi x}}{\lambda} = \frac{\sigma_*^2}{6M_P^2} \frac{\sqrt{4\pi}}{\lambda}$$  \hspace{1cm} (12)

where $\lambda \equiv \lambda_{\nu}/\sqrt{x}$ and $\Gamma_d \equiv x m$ is the inflaton decay rate which is smaller than $m$ with $x < 1$. Notice that the form of Eqs. (11) and (12) are the same and $\lambda$ here is just a parameter used to describe the case when the inflaton decay rate is smaller than the curvaton mass.

In this paper, we assume the curvature perturbation is dominated by curvaton and the curvature perturbation from inflaton is negligible. This liberates the constraint of inflation model building and allows the scale of inflation to be much lower. The condition is

$$P_{\zeta_{\text{inf}}}^{1/2} = \frac{1}{2\pi} \frac{H_* M_P}{\sqrt{\epsilon_H}} \lesssim 5 \times 10^{-5},$$  \hspace{1cm} (13)
where $\zeta_{inf}$ is the curvature perturbation generated from inflaton. For a typical value of $\epsilon_H \sim 0.01$, this implies $H_* \ll 10^{-5} M_P$. Using Eq. (8) and Eq. (12) and imposing CMB normalization ($P_{\zeta_\sigma}^{1/2} \simeq 5 \times 10^{-5}$), we obtain

$$\lambda = 3.9 \times 10^3 \sigma_* H_*$$

(14)

For conventional curvaton scenario like our case, there is a lower bound for the Hubble parameter during inflation, $H_* \gtrsim 10^7 \text{GeV} = 4.17 \times 10^{-12} M_P$ [24]. Hence we are interested in the range

$$10^{-11} \lesssim H_*/M_P \lesssim 10^{-6}.$$  

(15)

From Eq. (10), Eq. (12) and Eq. (14), we have

$$f_{NL} = 2.68 \times 10^5 \frac{H_*}{\sigma_*}$$

(16)

In this paper, we are interested in large non-Gaussianity and will explore the parameter space for $10 \lesssim f_{NL} \lesssim 100$. Using Eq. (14)-(16), we can plot $\lambda$ versus $H_*/M_P$ as shown in Fig. (1).

![Figure 1: $\lambda$ versus $H_*/M_P$](image)

Since curvaton must decay after inflaton $\Gamma = \lambda^2 m/4\pi < \Gamma_d = x m$, we obtain $x > \lambda^2/4\pi$ and $\lambda = \lambda_o/\sqrt{x} < \sqrt{4\pi}$. A lower bound for $\Gamma$ is obtained from the requirement that curvaton should not disturb big bang nucleosynthesis (BBN), which introduces [25]

$$\Gamma > 4.5 \times 10^{-25} \text{GeV} = 1.88 \times 10^{-43} M_P.$$  

(17)

An upper bound for $\Gamma_d$ can be obtained from the gravitino bound of reheating temperature if we assume the decay products of inflaton are thermalized immediately after decay, then we have [17]

$$T_R = \sqrt{\frac{\Gamma_d M_P}{k_{T_R}}} < 10^8 \text{GeV} = 4.17 \times 10^{-11} M_P$$

(18)

where $k_{T_R} = (4\pi^3 g(T_R)/45)^{1/2}$ and $g(T_R)$ is the effective number of massless degrees of freedom in thermal equilibrium. We will consider $k_{T_R} \approx 20$, corresponding to MSSM with $g(T_R) \approx 200$, this implies

$$\sqrt{\Gamma_d} = \sqrt{\lambda_o m} < 1.86 \times 10^{-10} M_P^{1/2}.$$  

(19)

\footnote{If the curvaton is a Pseudo Nambu-Goldstone Boson (PNGB) with a symmetry-breaking phase transition during inflation or the curvaton mass increases suddenly at some moment after the end of inflation, it is possible to get $H_* \text{ as low as } 1 \text{ TeV}$ [26, 27].}
Combining Eq. (17) and (19), we obtain the constraint
\[ 1.88 \times 10^{-43} < \frac{\lambda^2 m}{4\pi} < x m < 3.46 \times 10^{-20}. \] (20)

From Eq. (19) we can see that in order to evade the gravitino bound, small \( x \) is preferred, which means late decay of the inflaton, however this will cause the Yukawa coupling \( \lambda_\nu \) to be suppressed from Fig. (1). In Eq. (20), the upper bound can be relaxed if the inflaton does not thermalized immediately after decay, or if there is other methods to evade gravitino problem, for example, a period of thermal inflation \([28, 29]\) after reheating.

### 3 Initial condition

There is no consensus in literature about what the most natural initial condition for the curvaton field is. For example, if non-renormalizable terms are not protected by any symmetry, the slow roll condition may fail for large value of the curvaton field. This provides an upper limit for the curvaton field value, then one may choose the nature value of the curvaton field to be of the order of its upper limit. Usually the picture is that in different patches of the universes separated by the horizon, the curvaton field may take on different values at horizon exit, therefore the curvature perturbation and the amount of non-Gaussianities are different \([30, 31]\).

In \([25, 32]\), it is suggested that the most likely value of the curvaton field may be determined by the boundary between classical slow roll motion domination and quantum fluctuation domination. In one Hubble time, the classical slow roll gives a change \( \Delta \sigma = -V'/H_*^2 \), while the quantum fluctuation gives a random contribution \( \Delta \sigma = \pm H_*/2\pi \). When these two are equal, \( |V'| = H_*^3 \) and so gives \( \sigma_* \sim H_*^3/m^2 \). This suggests a value for the curvaton mass
\[ m^2 \sim \frac{H_*^3}{\sigma_*}. \] (21)

We will refer this case as case 1.

Another different argument is given in \([33]\) where the author gave three different arguments to suggest that the typical value of curvaton field is \( \sigma_* \sim H_*^2/m \). This suggests a curvaton mass
\[ m \sim \frac{H_*^2}{\sigma_*}. \] (22)

We will refer this case as case 2. We plot both cases in Fig. (2).

We should emphasis that those are not strict constraints to the allowed curvaton mass, however, the right-handed sneutrino mass may lie in the range where curvaton works best.

### 4 The spectral index

The spectral index \( n_s \) in curvaton scenario takes on the form
\[ n_s = 1 + 2\eta_{\sigma\sigma} - 2\epsilon, \] (23)
where
\[ \eta_{\sigma\sigma} \equiv \frac{1}{3H^2} \frac{d^2V(\sigma)}{d\sigma^2} \quad \text{and} \quad \epsilon \equiv -\frac{\dot{H}}{H^2}. \] (24)
Here $\eta_{\sigma \sigma}$ should be evaluated at horizon exit. WMAP data prefers a red tilted spectrum with spectral index $n_s \simeq 0.96$ [$14$]. Because a large positive $\eta_{\sigma \sigma}$ will result in a blue spectrum, it is not preferred. It is possible to get a large negative $\eta_{\sigma \sigma}$ during inflation. This happens for example from F-term hybrid inflation in which the vacuum energy during inflation breaks supersymmetry and will introduce a soft mass term of the order of the Hubble parameter. By choosing the magnitude correctly $2$, $n_s \simeq 0.96$ can be achieved [$17$, $34$]. We can imagine this also happens in D-term hybrid inflation if we consider a non-minimal gauge kinetic function, because a similar large mass correction can occur $35$ $36$. Another possibility is that if we have $\epsilon \sim 0.02$, $n_s \simeq 0.96$ can be achieved. However, this may not be easy to achieve for some models, for example, $[22]$. In this case, we will have $n_s \simeq 1$. In $[37]$, the authors argued that we can put $n_s = 1$ as the prior and show that it is not ruled out by WMAP data. Another way out is to assume there are some cosmic strings produced after inflaton as done in $[22]$.

5 Curvaton on the brane

As we can see from Eq. (20), it is generically true that for a successful curvaton model, the Yukawa coupling is very small. A very small Yukawa coupling may be regarded as fine-tuning. In this section, we will show that if we consider our world as a brane where gravity and matter field are confined on the 3-brane in the 5-dimensional spacetime $[38]$, the model can work with a larger Yukawa coupling. In this set-up, the Friedmann equation becomes $[39]$

$$H^2 \equiv \left( \frac{\dot{a}}{a} \right)^2 = \frac{\rho}{3M_P^2} \left( 1 + \frac{\rho}{2 \Lambda} \right),$$ (25)

where the brane tension $\Lambda$ relates the four dimensional Planck mass ($M_4 = 1.2 \times 10^{19}$ GeV) to the five dimensional Planck mass as $\Lambda = 3M_5^6/4\pi M_4^2$. The Big Bang Nucleosynthesis (BBN) limit implies $M_5 \gtrsim 10$ TeV $\sim 10^{-14}M_P$, where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. If $\rho/2\Lambda > 1$, we have $\rho \simeq \sqrt{6\Lambda H}M_P$, therefore $\rho_R(t_o) = \sqrt{6\Lambda}m_M$ and $\rho_R(t_D) = \sqrt{6\Lambda} \Gamma M_P$. Repeat

\footnote{This may be justified if we allow the curvaton mass to run via the method of $[40, 41]$}
what we did in Sec. 2, we found

$$\Omega_{\sigma,D} = \frac{m\sigma^2}{2\sqrt{6}\Lambda M_P} \left(\frac{m}{\Gamma}\right)^{1/4} = \frac{m\sigma^2}{2\sqrt{6}\Lambda M_P} \left(\frac{4\pi}{\lambda^2}\right)^{1/4}$$

(26)

and

$$P_{\sigma}^{1/2} = \frac{m\sigma}{6\sqrt{6}\Lambda M_P} \left(\frac{m}{\Gamma}\right)^{1/4} H_* = \frac{m\sigma}{6\sqrt{6}\Lambda M_P} \left(\frac{4\pi}{\lambda^2}\right)^{1/4} H_*.$$  

(27)

Here we have assumed $\rho/2\Lambda \lesssim 1$ before curvaton decay which implies

$$\Lambda \lesssim \frac{3}{2} \Gamma^2 M_P^2.$$  

(28)

For estimate, we saturate the inequality and by using Eqs. (26) and (27) we obtain

$$\lambda^{5/2} = 2.63 \times 10^4 \frac{\sigma_* H_*}{M_P^2}.$$  

(29)

For comparison, if $f_{NL} = 10$ (this implies $\sigma_* = 8.3 \times 10^2 H_*$), hence

$$\lambda^{5/2} = 6.58 \times 10^7 H_*^2 / M_P^2.$$  

(30)

Therefore, for example, when $H_* = 10^{-8} M_P$, we obtain $\lambda = 5.34 \times 10^{-4}$, which is larger than the usual case.

6 Conclusions

In this paper we explored the allowed value of Yukawa coupling $\lambda_\nu$ in the right-handed sneutrino curvaton scenario to generate large non-Gaussianity ($10 \lesssim f_{NL} \lesssim 100$) within the scale of inflation $10^{-11} \lesssim H_* \lesssim 10^{-6}$, which exhausted the allowed range for a normal curvaton model. The Yukawa coupling $\lambda_\nu$ will be further suppressed if the inflaton decay rate is lower than curvaton mass which is favored if we want to evade the gravitino problem in the framework of supersymmetry. We also consider two different kinds of scenarios in literature considering the most likely value of the curvaton field and show that these scenarios suggest the mass of the RH sneutrino. The mass is not severely constrained by this scenario, but it maybe of use if in the future we can determine the RH neutrino mass in a different way. We have also showed that in the case of braneworld scenario, the Yukawa coupling can be larger which may make it a more natural model.

In this paper our arguments are focused on the right-handed sneutrino curvaton, because it is a good candidate in the framework of MSSM. However, our analysis can also apply to the case of any curvaton models with a quadratic potential $V(\sigma) = m^2/2$ and decay rate of the form $\Gamma = h^2 m/4\pi$ where $h$ is some constant corresponds to our $\lambda_\nu$. In the models where supersymmetry is not imposed, there will be no gravitino bound for reheating temperature and $h$ is less constrained.

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