We study chaotic and non-oscillatory inflationary models in the curvaton scenario. In particular, we address the issue of large tensor-to-scalar ratio and red-tilted spectrum in chaotic models and reheating in non-oscillatory models in the light of latest Planck results. We show that curvaton can easily circumvent these problems and is well applicable to both type of models. For chaotic models, the observable non-Gaussianity put strong constraints on the decay epoch of curvaton and its field value around the horizon exit. Besides, it can also explain the observed red-tilt in the spectrum. As for the non-oscillatory models, curvaton by sudden decay into the background radiation provides an efficient reheating mechanism. To this effect, we consider the generalized exponential potential and obtain the allowed parametric space for model parameters. We then constrained the reheating temperature for both dominating and sub-dominating case, which satisfy the nucleosynthesis constraint.

I. INTRODUCTION

The latest Cosmic Microwave Background (CMB) observations [1, 3, 4] has outstandingly constrained non-Gaussianity, spectral index as well as tensor-to-scalar ratio of primordial perturbations. It is assumed that these primordial perturbations are an artifact of quantum fluctuations of all fields present during the so called ‘Inflationary era’. The viable scenario of inflation can be obtained either by a single scalar field (inflaton) [5–7] which sources both the inflation as well as primordial perturbations or by more than one scalar field [8–11] in which inflaton drives inflation whereas the subdominant field (curvaton) generates most of the primordial curvature perturbations [12–14].

In chaotic models, generally, both inflaton and curvaton contributes to overall density perturbations. When inflation ends, inflaton quickly decays into relativistic degrees of freedom and curvaton being a light scalar field initially, tends to become massive. This happens due to the fact that energy density of curvaton starts to increase as compared to that of the background radiation. The oscillating curvaton then also decays into radiation. If curvaton decays before dominating the total energy density of the Universe then it will lead to a significant non-gaussianity in the power spectrum which is disfavoured by latest CMB observations. Therefore one expects curvaton to decay after it becomes dominant species [15–17]. One of the main essence of taking into account another massive field like curvaton is to reduce the tensor-to-scalar ratio $r_T$ within the upper limit imposed by the Planck 2018 in order to save simple chaotic inflationary models from being ruled out [19, 20]. Moreover, the presence of curvaton can also explain the slight red-tilt in the scalar power spectrum as a consequence of its negative mass-squared value [33].

In non-oscillatory (NO) models, the inflaton field does not decay to background radiation but instead keeps rolling along its runaway type steep potential [28]. As a result, after a short while, its kinetic energy dominates the potential energy and the Universe enters into the so called ‘kinetic regime’. Since, the inflaton field potential does not have minimum, the standard reheating mechanism can not be applied here. Therefore, an alternative but efficient reheating mechanisms have been proposed [21, 23, 39] in which reheating is accomplished by the curvaton field based upon the decay of curvaton field. This mechanism has several
advantages over other mechanisms like perturbative decay, non-perturbative broad parametric resonance etc. as it provides sufficiently high reheating temperature for the standard nucleosynthesis process to occur. Therefore, due to its versatile behaviour, curvaton mechanism is well suited to both chaotic as well as non-oscillatory models.

The layout of the paper is as follows: In section, we study mixed inflaton-curvaton perturbations, and explore its effects in general chaotic inflationary models. For estimations, we particularly consider quadratic and quartic potentials. In section, we take generalized exponential potentials. In section, we take generalized exponential potentials and analyze the reheating mechanism by curvaton. In both cases, we consider quadratic potential for curvaton. For numerical estimations, we take \( N = 55 \) whenever required.

II. CHAOTIC MODELS: MIXED PERTURBATIONS

We begin with the general chaotic potential for the slow-roll inflaton field as

\[
V(\phi) = \frac{k \phi^\alpha}{\alpha M_{\text{pl}}^{\alpha-4}},
\]

where \( \alpha, k \) are dimensionless constants and \( M_{\text{pl}} \) is the Planck mass. In general, the slow-roll parameters are defined as

\[
\epsilon \equiv \frac{M_{\text{pl}}^2}{2} \left( \frac{1}{V} \frac{dV}{d\phi} \right)^2, \quad \eta_\phi \equiv \frac{M_{\text{pl}}^2}{V} \left( \frac{1}{V} \frac{d^2V}{d\phi^2} \right),
\]

which satisfy the condition \( \epsilon, |\eta_\phi| \ll 1 \) during inflation. Therefore, at the end of inflation, we have

\[
\epsilon_{\text{end}} = \frac{M_{\text{pl}}^2}{2} \left( \frac{\alpha}{\phi_{\text{end}}} \right)^2 = 1 \quad \Rightarrow \quad \phi_{\text{end}} = \frac{\alpha M_{\text{pl}}}{\sqrt{2}}.
\]

From the above expression, the duration of inflation or the number of e-folds \( N \) can be worked out in the following way:

\[
N \simeq \frac{1}{M_{\text{pl}}^2} \int_{\phi_{\text{end}}}^{\phi} \frac{V(\phi)}{V(\phi)} d\phi \simeq \frac{\phi^2}{2\alpha M_{\text{pl}}^2} - \frac{\alpha}{4},
\]

such that \( \epsilon \) and \( \eta_\phi \) can be expressed in terms of \( N \) as

\[
\epsilon = \frac{\alpha}{4N + \alpha}, \quad \eta_\phi = \frac{2(\alpha - 1)}{4N + \alpha}.
\]

Now, we assume a simple quadratic form of the potential for curvaton as

\[
V(\sigma) = \frac{1}{2} m_\sigma^2 \sigma^2,
\]

where \( \sigma \) and \( m_\sigma \) represents the curvaton field and mass, respectively. During inflation, curvaton ceases to be massless \( m_\sigma \ll H_{\text{end}} \) (where \( H_{\text{end}} := H \) at the end of inflation) and behaves as a light scalar field. But at the end of inflation, the parameter \( H \) starts to decrease as a consequence, curvaton eventually becomes massive \( m_\sigma \simeq H \). Curvaton then starts to oscillate about its mean field value in the early stages of radiation era and behaves as a pressureless matter component. While oscillating, its energy density \( \rho_\sigma \) varies as \( a^{-3} \) but that of the background radiation \( \rho_\gamma \) varies as \( a^{-4} \), where \( a \) is the scale factor. Therefore, after some time-intervals, curvaton starts to dominate the total energy density of the Universe.

The quantum fluctuations of curvaton after the horizon exit converts into primordial perturbations \( \xi \) (defined on the spatial slice of constant energy density). This is due to the fact that the size of these quantum fluctuations gets enhanced during inflation to the super-horizon scale and enters into the classical regime. These perturbations then remain frozen until they again enter inside the horizon which sets the initial conditions for perturbations i.e. \( \xi \simeq 10^{-5} \). Apart from adiabatic perturbations (also known as curvature perturbations), there might be some isocurvature perturbations present due to the difference between relative number densities of different components. Therefore, only the adiabatic perturbations are relevant for the structure formation. For weakly interacting inflaton and curvaton fields, one can write the resulting power spectrum \( P_\xi \) as

\[
P_\xi = P_\xi^\phi + P_\xi^\sigma = (1 + \lambda) P_\xi^\phi,
\]

where \( \lambda \equiv P_\xi^\sigma / P_\xi^\phi \) is the ratio of curvaton to inflaton power spectrum. By definition, the individual power spectrum of both fields are given as

\[
P_\xi^\phi = \frac{H_*^2}{8\pi^2 \xi M_{\text{pl}}^2}, \quad P_\xi^\sigma = \frac{r_d^2 H_*^2}{9\pi^2 \sigma_\xi^2},
\]

where \( H_* \) is the Hubble parameter, \( \sigma_* \) is the curvaton field evaluated just before the horizon exit and \( r_d \equiv \rho_\sigma / \rho \in [0, 1] \) is the ratio of curvaton energy density \( \rho_\sigma \) to the total energy density \( \rho \) of the Universe at the time of curvaton decay.

Now from Eqs. (7) and (8), we can express \( \lambda \) as

\[
\lambda = \frac{8}{9} \left( \frac{M_{\text{pl}}}{\sigma_*} \right)^2 \left( \frac{r_d}{r_d} \right)^2.
\]

It is clear from the above equation that for given initial conditions \( \epsilon \) and \( \sigma_* \), the extent of curvaton contribution
in the overall power spectrum is determined by \( r_d \). If it decays too early, its contribution in density perturbations will be small or vice-versa.

In the presence of curvaton spectral index \( n_s \) and tensor-to-scalar ratio \( r_T \) [20, 24] can be written as

\[
\begin{align*}
    n_s - 1 & = \frac{d \ln P_t}{d \ln k} = -2\epsilon + 2\eta_\sigma - \frac{4\epsilon - 2\eta_\phi}{1 + \lambda}, \\
    r_T & = \frac{P_T}{P_\xi^2 + P_\sigma^2} = \frac{16\epsilon}{1 + \lambda}, \tag{11}
\end{align*}
\]

where \( P_T \) is the tensor power spectrum and \( \eta_\sigma \) is the curvaton field slow-roll parameter similar to \( \eta_\phi \). So far, we have only considered the first-order of perturbations which has, of course, the major contribution in the spectrum. But CMB observations also encounters small but non-vanishing non-Gaussianity in the spectrum, mainly arising from the quadratic term. To take its effect into account, we write the non-Gaussianity parameter \( f_{NL} \) as [20]

\[
\frac{6}{5} f_{NL} = \left( \frac{1}{1 + \lambda} \right)^2 \frac{1}{2\epsilon} \left( 1 - \frac{\eta_\phi}{2\epsilon} \right) + \lambda^2 \left( 3 - 4r_d - 2r_d^2 \right). \tag{12}
\]

We now constrain these parameters in the light of latest Planck 2018 results.

**A. Planck Observational Constraints**

As we have mentioned in section [1] that a significant contribution of curvaton in overall density perturbations, to a certain extent, can alleviate the problem of large tensor-to-scalar ratio. To realize this, let us consider that if in Eq. (11), we impose the limit \( \lambda \rightarrow 0 \), then \( r_T \) will approach to 0.14 and 0.28 for \( \alpha = 2 \) and 4, respectively. Hence, theoretical estimates of \( r_T \) is too large to satisfy the upper-bound imposed by the joint Planck TT,TE,EE+lowE+lensing+BAO14 results i.e. \( r_T < 0.07 \) [3]. Due to this, the single field chaotic models are ruled out. Therefore, one finds that even a minor contribution of curvaton or a small non-zero value of \( \lambda \) can save these models by reducing \( r_T \). Therefore, in order to be consistent with observations, we obtain the lower bound on \( \lambda \) from Eq. (11) as

\[
\lambda > \frac{228.5\alpha}{4N + \alpha} - 1. \tag{13}
\]

Also from Eqs. (13), (5) and (10), it follows that

\[
n_s - 1 > \frac{-2\alpha}{4N + \alpha} + 2\eta_\sigma - \frac{4}{228.5\alpha} \tag{14}
\]

which implies

\[
\eta_\sigma < \frac{1}{2} \left( \frac{0.0175}{\alpha} - 1 + n_s + \frac{2\alpha}{4N + \alpha} \right). \tag{15}
\]

By using the above expression, we have depicted the allowed region between \( \eta_\sigma \) and \( n_s \) for quadratic and quartic potential (see fig. (1)). In the figure, it can be seen that the observational estimates of \( n_s \) restrain \( \eta_\sigma \) to always remain negative for \( \alpha = 2 \), whereas a slight positive value is still allowed for \( \alpha = 4 \). Now, in order to obtain theoretical upper bound on \( \eta_\sigma \), we use Planck TT,TE,EE+lowE+lensing+BAO 2018 best-fit \( n_s = 0.9665 \pm 0.0038 \), which gives \( \eta_\sigma < -0.003 \) for \( \alpha = 2 \) and \( \eta_\sigma < 0.003 \) for \( \alpha = 4 \).

Therefore, in order to explain observed red-tilted spectrum \( n_s < 1 \) which is 9\( \sigma \) level away from scale-invariance \( \eta_\sigma \), curvaton needs to have negative \( \eta_\sigma \). However, the opposite is not true as one can also get flat spectrum \( (n_s = 1) \) with negative \( \eta_\sigma \) (see fig. (1)).

Now, by using Eq. (15), the expression (12) can be expressed in terms of \( N \) and \( \alpha \) as

\[
\frac{6}{5} f_{NL} = \left( \frac{1}{1 + \lambda} \right)^2 \frac{2}{4N + \alpha} + \lambda^2 \left( 3 - 4r_d - 2r_d^2 \right). \tag{16}
\]

In mixed perturbations, if contribution of inflaton is comparable to or greater than that of the curvaton, then it may lead to large non-Gaussianity in the spectrum, a signature of which is clearly absent in the Planck results. Hence, one expects curvaton to decay only when it gets dominated. The observable non-Gaussianity puts strong limit on the

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**FIG. 1:** The parametric space between \( n_s \) and \( \eta_\sigma \) for \( \alpha = 2 \) and 4. The horizontal dotted lines represents corresponds to Planck results \( n_s = 0.9665 \pm 0.0038 \).
decay epoch of curvaton and its field value $\sigma_*$. In particular, in Fig. (2), we have depicted the corresponding dependence between $r_d$ and $\sigma_*$ by using above equation [16] for different level of non-Gaussianity. In same figure, one can observe that $f_{NL}$ increases as $\sigma_*/M_{pl}$ decreases, whereas it increases with $r_d$. Hence, one can impose the observational bounds on $f_{NL}$, on $\sigma_*$. In order to do this, we use the observational results $f_{NL} = 22.7 \pm 25.5$ [2, 35] and obtain

$$10^{-4} M_{pl} \leq \sigma_* \leq 10^{-1} M_{pl}, \quad \text{with} \quad 0.016 \leq r_d \leq 1, \quad (17)$$

which supports the fact that curvaton should decay after become dominant and the field $\sigma_*$ should remain sub-Planckian near the horizon exit.

### III. NO MODEL: GENERALIZED EXPONENTIAL POTENTIAL

In this section, we shall consider a class of *quintessential inflationary* models which is based on the assumption that a single inflaton field can source both early and late-time cosmic acceleration [28, 30, 36, 37]. In these models, the field potential is of simple exponential type, at least in the post inflationary era, a generalization of which is given as

$$V(\phi) = k \exp \left[ -\lambda \left( \frac{\phi}{M_{pl}} \right)^n \right] M_{pl}^4, \quad (18)$$

where $\lambda$ and $n$ are constants. The potential $V(\phi)$ has an interesting behaviour that during inflation it remains shallow but becomes steep in the post inflationary era \(^1\). Also at late times, it gives rise to an approximate scaling solution as $\Gamma \equiv V_{\phi\phi}V/\phi^2 \rightarrow 1$ for large $\phi$ value which, in order to exit from the scaling regime and to source late-time cosmic acceleration, needs to be non-minimally coupled to matter (for more details, see ref. [38, 42]).

The standard slow-roll parameters for this model are given as

$$\epsilon = \frac{1}{2} n^2 \lambda^2 \left( \frac{\phi}{M_{pl}} \right)^{2n-2}, \quad (19)$$

$$\eta_{\phi} = -n \lambda \left[ n - 1 - n \lambda \left( \frac{\phi}{M_{pl}} \right)^n \right] \left( \frac{\phi}{M_{pl}} \right)^{n-2}, \quad (20)$$

such that the violation of the slow-roll condition $\epsilon|_{\phi=\phi_{end}} = 1$ confirms the end of the inflationary period. As a result, one can estimate field at the end of inflation i.e. $\phi_{end}$ as

$$\phi_{end} = \left( \frac{2}{n^2 \lambda^2} \right)^{\frac{1}{2n-2}} M_{pl}, \quad (21)$$

which in order to behave like a quintessence field in late-times should satisfy $\phi \gg M_{pl}$ condition, and that can only happen if $\lambda \ll 1$. Also, we obtain the number of e-folds $N$ as

$$N = \frac{1}{n \lambda (n-2)} \left[ \left( \frac{\phi}{M_{pl}} \right)^{2-n} - \left( \frac{2}{n^2 \lambda^2} \right)^{\frac{2-n}{2n-2}} \right] \left( \frac{2}{n^2 \lambda^2} \right)^{\frac{1}{2n-2}}, \quad (22)$$

which is valid for any $n > 1$ except $n = 2$, as there exist a singularity. Now, by re-expressing the above equation in terms of $\phi$ as

$$\phi = \left[ n(n-2)\lambda N + \left( \frac{2}{n^2 \lambda^2} \right)^{\frac{2-n}{2n-2}} \right] \left( \frac{2}{n^2 \lambda^2} \right)^{\frac{1}{2n-2}} M_{pl}, \quad (23)$$

we obtain a simplified expression by imposing the large-field limit for the reason mentioned earlier

$$\phi \approx \left[ n(n-2)\lambda N \right]^{\frac{1}{2n-2}} M_{pl}. \quad (24)$$

By using the constraint $r_T = 16 \epsilon < 0.07$ [3], we obtain the parametric space between $n$ and $\lambda$ (see Fig. (3)) in which the shaded region represents the allowed parametric space whereas the white portion is excluded. Since, we have

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\(^1\) In this setup one needs to shift the field $\phi$ which is not justified in the absence of shift symmetry.
already stated before that the large-field limit demands \( \lambda \ll 1 \), it is clear from Fig. (3) that to satisfy this condition \( n \) must be greater than unity (except \( n \neq 2 \)). So in order to give rise to quintessential effects at late times, exponential potential with \( n > 1 \) is favoured over \( n = 1 \).

Now, to estimate \( \lambda \) for each \( n \), let us consider the standard expression of spectral index \( n_s = 1 - 6 \epsilon + 2\eta_{\phi} \), which can be written in more explicit form by using Eq. (19) and (20) as

\[
\begin{aligned}
\quad n_s - 1 &= -n\lambda \left[ 2(n - 1) + n\lambda \left( \frac{\phi}{M_{pl}} \right)^n \right] \left( \frac{\phi}{M_{pl}} \right)^{n-2}.
\end{aligned}
\]

By again making use of the best fit value \( n_s = 0.9665 \) (Planck TT,TE,EE+lowE+lensing+BAO 2018), we obtain \( \lambda = 5.6 \times 10^{-6}, 2.8 \times 10^{-10} \) and \( 7.8 \times 10^{-15} \) for \( n = 4, 6 \) and 8, respectively. Also, by considering the COBE normalization \( 41 \) i.e. \( k^{1/4} = 0.013 r_\gamma^{1/4} \) together with Eqs. (19) and (23), we obtain

\[
\begin{aligned}
\quad k &\approx 2.8 \times 10^{-8} \left( n^2 \lambda^2 [n\lambda(n-2)] \times \frac{2(n-1)}{n-1} \right) \\
&\approx 9.5 \times 10^{-10}.
\end{aligned}
\]

for all obtained sets of \( n \) and \( \lambda \). By using these estimations, we will now constrain the reheating temperature.

As we know that in NO models, the inflaton field does not decay and therefore an alternative source to execute the reheating mechanism is required. For this purpose, one requires another light scalar field like curvaton \( m_\sigma \ll H_{inf} \), which can take care of reheating process. After inflation ends, the Universe enters into the kinetic regime and curvaton starts to oscillate about its mean field value and finally becomes massive \( m_\sigma \sim H_{end} \). But in order to prevent another inflationary scenario, curvaton still remains sub-dominant in the beginning of kinetic regime, by satisfying the following condition

\[
\rho_\sigma \ll \rho_\phi = 3H_{end}^2 M_{pl}^2 \quad \Rightarrow \quad \sigma_\sigma^2 \ll \frac{3}{4\pi} M_{pl}^2, \quad (27)
\]

where \( \rho_\sigma \approx m^2 \sigma_\sigma^2 \). Here, we have assumed that \( \sigma_\sigma \approx \sigma_i \) (where \( \sigma_i \) is the initial field value). Also, the sub-dominant condition of curvaton during inflation constraints the curvaton mass as

\[
\begin{aligned}
\quad V(\sigma) &= \frac{1}{2} k \exp[-\lambda (\frac{\phi}{M_{pl}})^n] M_{pl}^4 \\
&\approx m_\sigma \left( \frac{8\pi k \exp[-\lambda (\frac{\phi}{M_{pl}})^n]}{3} \right)^{1/2} M_{pl}, \quad (28)
\end{aligned}
\]

where we have used Eq. (27). Now for the obtained values of \( n, \lambda \) and \( k \), we obtain the upper bound on \( m_\sigma \) as

\[
\quad m_\sigma \ll 2.89 \times 10^{-5} M_{pl}, \quad (29)
\]

which fulfills the above said requirement \( m_\sigma \ll H_{inf} \), as \( H_{inf} \lesssim 10^{-6} \) during inflation.

As we have stated above that in NO models curvaton by sudden decay, creates all the matter present in the Universe. Therefore, by using the standard definition of decay parameter \( \Gamma_\sigma \), we can constrain its decay epoch. For dominating case when \( \rho_\sigma > \rho_\phi \), \( \Gamma_\sigma \) satisfies the following condition \( 21 \)

\[
\quad \frac{\Gamma_\sigma}{m_\sigma} \leq \frac{4\pi \sigma_\sigma^2}{3M_{pl}^2} < 1 \quad (30)
\]

and its corresponding reheating temperature \( T_{rh} \sim \rho_\sigma^{1/4} \sim \sqrt{3M_{pl} \Gamma_\sigma} \) \( 36 \) satisfies

\[
\quad T_{rh} \lesssim 8.51 \times 10^{-9} M_{pl}. \quad (31)
\]

where, by keeping in mind above constraints (27) and (28), we assume \( m_\sigma \approx 10^{-7} M_{pl} \) and \( \sigma_\sigma \approx 10^{-5} M_{pl} \).

Similarly, for the sub-dominating case when \( \rho_\sigma < \rho_\phi \), \( \Gamma_\sigma \) satisfies

\[
\quad \frac{4\pi m_\sigma \sigma_\sigma^2}{3 M_{pl}^2} \leq \Gamma_\sigma \leq m_\sigma, \quad (32)
\]

and \( T_{rh} \) is given as

\[
\quad T_{rh} \approx \sqrt{\frac{m_\sigma^3 \sigma_\sigma^3}{3 M_{pl}^4 \Gamma_\sigma^{1/2}}}. \quad (33)
\]
which after rearranging and plugging back in Eq. (32), we get
\[ \sqrt{\frac{m_\sigma \sigma_1^2}{3M_{pl}^2}} \leq T_{re} \leq \sqrt{\frac{m_\sigma \sigma_2^2}{3M_{pl}^2}}, \]
which implies
\[ 5.77 \times 10^{-12} M_{pl} \leq T_{re} \leq 2.40 \times 10^{-9} M_{pl}, \]
which is, as expected, well satisfies the requirement for the standard big bang nucleosynthesis (BBN) process to occur. Although, one can satisfy BBN constraint for a wide range of model parameters \( n \) and \( \lambda \), it is still interesting to note that curvaton reheating indirectly depends on the inflaton field which do not take part in the reheating process.

IV. DISCUSSION AND CONCLUSIONS

In this paper, we have examined the viability of both chaotic as well as non-oscillatory inflationary models. For chaotic models, we have carried out analysis for two types of potential namely, quadratic and quartic. We have shown that for both forms of potentials the presence of curvaton field can indeed alleviate the problem of having large tensor-to-scalar ratio specific to the single field inflationary models. Also from the observed red-tilt in the spectrum, we have obtained upper bound on \( n_g \) and show that for quadratic potential it is always negative but for quartic potential it can also take small positive value. Moreover, the observational bound on the non-gaussianity parameter constraints curvaton to decay after it gets dominated.

As for the NO models, which are characterized by a run away type potential, the inflaton field survives to account for late time physics. We have thus considered the generalized exponential potential which can successfully account for inflation. After inflation, the field potential becomes steep and despite the fact it is not exponential, it might give rise to scaling behaviour in the asymptotic regime as \( \Gamma \equiv V_\phi V_{\phi^2} \rightarrow 1 \) for large values of the field. In this case, one could use an alternative reheating mechanism based on the curvaton decay, which interacts with inflaton only gravitationally. In this paper, we have explored that curvaton reheating seems to be an ideal in this case.

As for the parameter estimation, we have again considered Planck 2018 results and have depicted the parametric space between \( n \) and \( \lambda \) and obtain their possible set of values. We have estimated parameter \( k \approx 9.5 \times 10^{-10} \) and correspondingly allowed limits for \( T_{re} \) for dominating as well as sub-dominating case.

We have thus demonstrated that curvaton scenario is appropriate to both chaotic as well as NO models, in particular, quintessential inflationary models.

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