The inflaton-curvaton scenario in the MSSM and predictions for non-Gaussianity

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We provide a model in which both the inflaton and the curvaton are obtained from within the minimal supersymmetric Standard Model, with known gauge and Yukawa interactions. Since now both the inflaton and curvaton fields are successfully embedded within a visible sector, their decay products thermalize very quickly and before the electroweak scale. This results in two important features of the model: firstly, there will be no residual isocurvature perturbations, and secondly, observable non-Gaussianities may be generated with the non-Gaussianity parameter $f_{NL} \sim \mathcal{O}(10-1000)$ being determined solely by microscopic parameters which can be constrained at the LHC.

The curvaton scenario [1–3] is an alternative mechanism for the generation of the primordial perturbations whose spectrum is observed in the cosmic microwave background (CMB) [4]. In this scenario, the density perturbations are sourced by the quantum fluctuations of a light scalar field ϕ, the curvaton, which makes a negligible contribution to the energy density during inflation and decays after the decay of the inflaton σ into radiation. (For a review on inflation and the curvaton mechanism, see [5].)

If the curvaton dominates energy density at the time of its decay, it would be solely responsible for creating the source perturbations for the CMB anisotropies, and exciting all the Standard Model (SM) degrees of freedom (*dof*) [6, 7]. However, if the curvaton does not totally dominate while decaying, then it might leave two potentially observable imprints:

Isocurvature perturbations: If either the inflaton or curvaton belong to a hidden sector, then they may couple to a myriad other hidden sectors beyond the SM as well to the SM *dof*. There is no guarantee that the hidden and visible sector *dof* should reach thermal equilibrium before Big Bang Nucleosynthesis (BBN) [8] takes place. In such a case, residual isocurvature perturbations are expected. These are tightly constrained by the CMB data to be less than 10% [4].

Primordial non-Gaussianity: The initial isocurvature perturbations of the curvaton field are converted into adiabatic curvature perturbations, and as the curvaton oscillates or rotates, non-Gaussian fluctuations can be enhanced to the level which can be constrained by the PLANCK mission. The enhancement in non-Gaussianity is typically given by $f_{NL} \sim 5/(4r)$ for $r < 1$, where $r \equiv \rho_\phi/\rho_{\rm rad}$ at the time the curvaton decays [1]. The factor r is also known as the inefficiency factor.

The challenge for any *viable* curvaton scenario in which the curvaton leaves its imprints on non-Gaussianity is two-fold: firstly the curvaton must be light during inflation, and secondly the inflaton decay products and the curvaton decay products must thermalize before the time of nucleosynthesis, as there are stringent constraints on any non-SM like hidden radiation after BBN [8]. In order to achieve this, we wish to the entire inflation-curvaton paradigm within visible-sector physics, which is a daunting task in itself, given that only very recently we have understood how to embed the inflationary paradigm within a visible sector with known gauge interactions [9].

The aim of this paper is to show, for the first time, that it is indeed possible to embed the inflation and curvaton paradigms within a viable visible sector model beyond the SM, where all the interactions are solely governed by the SM gauge interactions, without involving any hidden sector. We will do this in the context of the MSSM (minimal supersymmetric Standard Model) by showing that there exists an interesting landscape constructed by *gauge invariant* field operators for which one can find both inflaton and curvaton candidates within the MSSM. We will thus provide a cosmological solution to a general problem of the curvaton scenario, i.e. how to generate measurable non-Gaussianity without large residual isocurvature fluctuations.

Let us first consider the total potential to be the sum of inflaton vacuum, denoted by V_0, and curvaton potential $V(\phi)$

$$V_{total} = V_0 + V(\phi)\,. \tag{1}$$

We assume $V''(\phi) \sim m_\phi^2(\phi_I) \ll H_I^2 \sim V_0/M_{\rm P}^2$ ($M_{\rm P} \sim 10^{18}$ GeV) where the subscript I indicates the quantities are evaluated during inflation. This condition is required for a successful curvaton scenario. The curvaton acquires vacuum induced quantum fluctuations, which have amplitude

$$\delta = \frac{H_I}{2\pi\phi_I}\,. \tag{2}$$

These fluctuations are converted into the adiabatic density perturbations when the curvaton decays during its coherent oscillations or rotations. In order to match the observed amplitude of the fluctuations on the CMB, $\delta \sim 10^{-5}$.

Let us first discuss the origin of the curvaton, which we take to be an R-parity conserving D-flat direction of the MSSM (for a review see [10]). The two candidate fields are LLe (where L denotes the left-handed slepton superfield and e the right-handed superfield), or udd (where u and d denote the right-handed squark superfields), which are lifted by the operator

$$W \supset \frac{\lambda}{n}\frac{\Phi^n}{M_{\rm P}^{n-3}}\,, \qquad (3)$$

where $\lambda \sim \mathcal{O}(1)$. The scalar component of the Φ superfield is denoted by ϕ. At the lowest order the potential along the ϕ direction is given by:

$$V(\phi) = \frac{m_\phi^2|\phi|^2}{2} + \lambda^2\frac{|\phi|^{2(n-1)}}{M_{\rm P}^{2n-3}} + \left(A\lambda\frac{\phi^n}{M_{\rm P}^{n-3}} + h.c.\right)\,, \qquad (4)$$

where $A \sim m_\phi \sim \mathcal{O}(100-1000)$ GeV, m_ϕ is the soft SUSY-breaking mass term, and $n = 6$ for udd, LLe. During inflation if $m_\phi^2 \ll H_I^2$, the fluctuations along this nearly massless direction would create a homogeneous condensate with a VEV given by

$$\phi_I \sim \left(m_\phi M_{\rm P}^{n-3}\right)^{1/n-2} \sim 10^{14}\ \text{GeV}. \qquad (5)$$

For $m_\phi \sim 100 - 1000$ GeV, and $n = 6$, in order to match the amplitude of the density perturbations δ, the Hubble expansion rate during inflation should be $H_I \sim 10^{10}$ GeV.

There is a distinction between a positive and a negative phase of the A term. The difference in dynamics arises after the end of inflation. In the case of positive A-term the curvaton starts rolling towars the origin immediately, but in the case of a negative phase, for values of $A \geq \sqrt{40}m_\phi$, it may remain in a false vacuum with the VEV given by Eq. (5). In this case the curvaton rotates instead of oscillates around its global minimum at $\phi = 0$. In either scenario, the curvaton mass is negligible compared to the Hubble expansion rate. In fact, for $A = \sqrt{40}m_\phi$ and a negative phase the curvaton is actually massless along the real direction, and obtains inflaton-induced random fluctuations of order $\delta\phi \approx H_I/2\pi$.

We now provide two distinct possibilities for the origin of V_0 within the MSSM.

(a) False vacuum from MSSM landscape: Let us consider a particular combination from the renormalizable part of the MSSM superpotential, QH_uu, the monomial with associated Yukawa coupling yQH_uu, where Q, u, and H_u denote the left-handed squark, right-handed up-type squark and the Higgs (which gives mass to the up-type quarks) fields. It represents a D-flat direction $\sigma = (\frac{1}{\sqrt{3}})(Q + H_u + u)$. The gauge invariance allows higher order superpotential terms, i.e. $(QH_uu)^n$ and $(QH_uu)^n(H_uH_d)$. After rewriting them in terms of σ and $\chi \equiv H_d$, the superpotential reads (to the lowest order) [11]

$$W \supset y\frac{\sigma^3}{3} + \lambda\frac{\sigma^6}{6M_{\rm P}^3} + \lambda'\frac{\sigma^4\chi}{4M_{\rm P}^2} + \ldots, \qquad (6)$$

where λ, λ' are constants of $\mathcal{O}(1)$. The scalar potential along the flat direction follows (note that here we denote both the superfield and the field by σ):

$$V(\sigma) = |y\sigma^2 + \lambda\frac{\sigma^5}{M_{\rm P}^3}|^2 + \left(\frac{\lambda'^2|\sigma|^8}{16M_{\rm P}^4}\right) + \ldots. \qquad (7)$$

We would also have contributions from soft SUSY-breaking terms, but these terms are subdominant at higher VEVs. The first term on the right-hand side of Eq. (7) vanishes at the origin $\sigma = 0$ and three other points with radial and angular coordinates (σ_0, θ_0) where $\sigma_0 = (y/\lambda)^{1/3}M_{\rm P}$, $3\theta_0 + \theta_\lambda - \theta_y = \pi,\ 3\pi,\ 5\pi$. Here θ_y, θ_λ are the phases of the y and λ couplings respectively. Note that $y \ll 1$ except for the top quark Yukawa, and hence $\sigma_0 \ll M_{\rm P}$. This implies that the higher order terms in Eq. (6) (denoted by ...) have a subleading contribution. These minima are separated from the origin by a barrier whose height is given by [11]

$$V_{\rm barrier} = \frac{9}{25}\left(\frac{2}{5}\right)^{4/3} y^2\left(\frac{y}{\lambda}\right)^{4/3} M_{\rm P}^4. \qquad (8)$$

The second term on the right-hand side of Eq. (7) is positive-definite and lifts the potential at $|\sigma| \neq 0$. This results in having three false minima at $|\sigma| \sim \sigma_0$. Depending on the relative size of y, λ and λ', the potential at these false minima can assume any value $V_0 \lesssim V_{\rm barrier}$. False vacuum inflation in these minima has a Hubble expansion rate

$$H_I \lesssim \frac{V_{\rm barrier}^{1/2}}{\sqrt{3}M_{\rm P}} \lesssim y^{5/3}M_{\rm P}\,. \qquad (9)$$

For $y \sim 10^{-5} - 10^{-2}$ (which is the case for all of the SM Yukawa couplings except for the top quark) the false vacuum inflation could be driven at a Hubble rate as large as $H_I \sim 10^{10} - 10^{14}$ GeV for $\lambda \sim \lambda' \sim \mathcal{O}(1)$.

If the inflaton is locked in a false vacuum inflation can happen eternally, but it would eventually tunnel to the true vacuum, $\sigma = 0$, by bubble nucleation. However, the curvaton already belongs to the true vacuum. The fluctuations of the curvaton ϕ would lead to the structure formation. In this case the ideal curvaton candidate would be: $\phi = \left(\frac{1}{\sqrt{3}}\right)(L + L + e)$ which could also be lifted simultaneously. Given a rapid tunnelling rate and bubble collisions, which is possible within the MSSM, the potential energy density V_0 can eventually be transferred to a thermal bath of MSSM *dof.*

(b) Slow roll inflation within MSSM landscape: In this scenario we consider two flat directions, LLe and udd, one of which is the inflaton and the other the curvaton. We take the inflaton direction to be σ, and gauge invariance here allows terms like $m = 2, 3, 4, \ldots$

$$W = \sum_m \frac{\lambda_m}{3m} \frac{\sigma^{3m}}{M_{\rm P}^{3m-3}} \,. \tag{10}$$

The potential at the lowest order would be:

$$V = \left| \lambda_2 \frac{\sigma^5}{M_{\rm P}^3} + \lambda_3 \frac{\sigma^8}{M_{\rm P}^6} + \lambda_4 \frac{\sigma^{11}}{M_{\rm P}^9} + \ldots \right|^2 \tag{11}$$

where $\ldots$ contain the higher order terms. Such potentials were studied in Refs. [12, 13]. For $\lambda_2 \ll \lambda_3 \ll \lambda_4 \ll \lambda_n \leq \mathcal{O}(1)$, such potentials provide a unique solution for which first and second derivatives of the potential vanish along both radial and angular direction in the complex plane: $\partial V/\partial\sigma = \partial V/\partial\sigma^* = \partial^2 V/\partial\sigma^2 = \partial^2 V/\partial\sigma^{*2} = 0$ (a saddle point condition) [14]. For the first three terms in Eq. (11), it is possible to show that this happens when

$$\lambda_3^2 = \frac{55}{16}\lambda_2\lambda_4 \,, \tag{12}$$

at the VEVs: $\sigma = \sigma_0 \exp[i\pi/3,\ i\pi,\ i5\pi/3]$, $\sigma_0 = (2/11)(\lambda_3/\lambda_4)^{1/3} M_{\rm P}$. Concentrating on the real direction, the potential energy density stored in the inflaton sector is given by:

$$V_0 \sim \left(\frac{153}{88}\right)^2 \lambda_2^2 \frac{\sigma_0^5}{M_{\rm P}} \,, \tag{13}$$

where $\sigma_0 \ll M_{\rm P}$. Note that inflation happens near the saddle point σ_0, where the effective mass vanishes. However, the third derivative of the potential is not negligible, $V''' \sim \lambda_2^2 \sigma_0^2/M_{\rm P} \neq 0$, which leads to the end of slow roll inflation. The corresponding Hubble expansion rate is given by $H_I \sim \lambda_2^2 \sigma_0^5/M_{\rm P}^3$. For $\sigma_0 \sim 10^{14}$ GeV and $\lambda_2 \sim 10^{-3} - 10^{-4}$, it is possible to obtain $H_I \sim 10^{10}$ GeV, required for a successful curvaton scenario.

Now let us consider the aftermath of inflation. In either of the cases (a) or (b), the inflaton would decay primarily into MSSM *dof.* The coherent oscillations of the inflaton would give rise to instant preheating and thermalization of the light MSSM *dof* as discussed in Ref. [15], with a reheat temperature

$$T_R \sim [H_I M_{\rm P}]^{1/2} \sim 10^{13}\ {\rm GeV} \,. \tag{14}$$

Note that depending on the nature of flat direction curvaton, not *all* of the MSSM *dof* need be in thermal equilibrium. For instance, consider the case where inflation is driven by either $QH_u u$ in scenario (a), or udd in scenario (b), and the curvaton direction is LLe. If both inflaton and curvaton simultaneously take large VEVs, the $SU(2)_W$ *dof* would not reach thermal equilibrium, since the LLe VEV would induce large masses to those *dof.* This can play a crucial role in determining the non-Gaussianity parameter f_{NL}, as we shall show below.

The curvaton ϕ starts to rotate about the origin when $H = H_{\rm osc} \sim m_\phi$. The field value at this time is $|\phi_{\rm osc}| \sim (m_\phi M_{\rm P}^{n-3})^{1/n-2}$. During this epoch the universe is already radiation-dominated following the decay of the inflaton. However, the curvaton cannot decay immediately, due to the fact that the curvaton VEV induces large masses $h\langle\phi(t)\rangle$ for gauge bosons, gauginos and (s)leptons, where h is the gauge or Yukawa coupling. The curvaton's decay at leading order is *kinematically* forbidden if $h\langle\phi\rangle \geq m_\phi/2$. Decays do not occur until the Hubble expansion has redshifted ϕ down to $m_\phi/2h$. Note that the Yukawa couplings are typically smaller than the gauge couplings. During the rotations, the curvaton VEV will scale as $\phi(t) \propto a^{-3/2}$, as $a \propto H^{-1/2}$ during radiation dominated epoch. Therefore, the curvaton decays when [16]

$$H = H_{\rm dec} \sim m_\phi \left(\frac{m_\phi}{h\phi(t)}\right)^{4/3} \,, \tag{15}$$

For large ϕ, the decay time is naturally longer than the normal decay rate into the massless *dof.* The radiation energy density stored in the inflaton decay products scales as $\rho_{vis} \propto H^2$, where the subscript denotes the visible *dof.* The ratio of the energy densities at the time the curvaton decays is given by

$$\begin{aligned} r \ \equiv\ \frac{\rho_\phi}{\rho_{\rm vis}} &\sim \left.\frac{\rho_\phi}{\rho_{\rm vis}}\right|_{\rm osc} \left(\frac{H_{\rm dec}}{H_{\rm osc}}\right)^{-1/2} \,, \\ &\sim \left(\frac{m_\phi}{M_{\rm P}}\right)^{2/(n-2)} \left(\frac{m_\phi}{h\phi}\right)^{-2/3} \leq 1 \,. \end{aligned} \tag{16}$$

The *kinematical blocking* due to the curvaton VEV enhances the inefficiency factor, r, therefore the curvaton rotations prolong the mater dominated epoch till it decays completely. For soft SUSY-breaking mass $m_\phi \lesssim 1000$ GeV, the inefficiency parameter is $r \sim \mathcal{O}(1)h^{2/3}$. Depending on the SM Yukawas involved in the interactions, the non-Gaussianity parameter would then be expected to be in the range

$$f_{NL} \sim \frac{5}{4r} \sim \mathcal{O}(1)h^{-2/3} \sim 10 - 10^3 \,, \tag{17}$$

for $h \sim 10^{-2} - 10^{-5}$. The temperature at which the curvaton decay products reach thermal equilibrium is determined by Eq. (15). A thermal bath filled with MSSM *dof* would be obtained by

$$T \sim (H_{\rm dec} M_{\rm P})^{1/2} \sim 10^{4.5} - 10^{6.5}\ GeV \tag{18}$$

for $h \sim 10^{-2} - 10^{-5}$. Such a temperature is sufficient to excite weakly interacting massive particles and for baryogenesis [17]. Note that both the temperatures from Eqs. (14) and (18) are sufficiently high to excite thermal/non-thermal gravitinos. If the gravitinos are the lightest SUSY particle, this causes two problems for this scenario: over-production of gravitinos with both helicities would be bad for BBN, and the gravitinos would thermally decouple even before the curvaton has started decaying. This would generate large residual isocurvature perturbations, because gravitinos can never come into thermal equilibrium. Instead the ideal dark matter candidate would be the neutralino, which decouples from the thermal plasma at $T \sim 40 - 50$ GeV.

To summarize, the MSSM landscape has interesting features—from the low energy point of view there exist regions of false vacua, saddle points, and so on—which allow us to construct a model in which both the inflaton and curvaton can be embedded within a visible sector of the MSSM. The radiation created from the decay of the inflaton and curvaton belong to the visible sector, avoiding the problem of residual isocurvature fluctuations. The curvaton mechanism in this model can create observable non-Gaussianity. The non-Gaussianity parameter f_{NL} ranges from $10 - 1000$, and depends crucially on the microscopic details of the particle physics, which will be tested directly at the LHC. Furthermore, the model favours a visible-sector dark matter candidate such as the lightest neutralino.

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