A three fluid system describing the decay of the curvaton is studied by numerical and analytical means. We place constraints on the allowed interaction strengths between the fluids and initial curvaton density by requiring that the curvaton decays before nucleosynthesis while nucleosynthesis, radiation-matter equality and decoupling occur at correct temperatures. We find that with a continuous, time-independent interaction, a small initial curvaton density is naturally preferred along with a low reheating temperature. Allowing for a time-dependent interaction, this constraint can be relaxed. In both cases, a purely adiabatic final state can be generated, but not without fine-tuning. Unlike in the two fluid system, the time-dependent interactions are found to have a small effect on the curvature perturbation itself due to the different nature of the system. The presence of non-gaussianity in the model is discussed.

I. INTRODUCTION

The problem of determining the evolution of large scale perturbations in a background of a multi-component fluid system is central in modern day cosmology [1, 2, 3]. In such a system, the interactions between the different fluids are important in determining the evolution of the curvature perturbation [4]. Examples of interacting fluid systems demonstrate the importance of such systems as most notably reheating at the end of inflation [10, 11, 13, 14, 16, 17] and the curvaton scenario [18, 19, 21, 26, 27].

In contrast to any single fluid system, in a multi-component system the total curvature perturbation, \( \zeta \), generally evolves whenever the non-adiabatic pressure is non-zero, i.e. when interactions between the fluids exist. Evolution of the primordial large scale curvature perturbation can relax the underlying assumptions on the inflationary scenario. Therefore analysis of multi-component fluid systems may affect our view on the physical settings. In addition to the curvature scenario considered in this paper, natural frameworks for such mechanism exist e.g. within a traditional multiple inflationary scenario [33] or a string landscape picture [34, 35]. Whether a given scenario can effectively modify the primordial spectrum depends on the exact nature of the system.

Recent cosmic microwave surveys have also brought attention into the concept of non-gaussianity i.e. how much the spectrum deviates from gaussian distribution. This is especially important in multifield models of inflation including the curvaton scenario [7, 9].

The mechanism how energy is transferred between the fluids can be described by different methods, e.g. by a constant interaction [5] or by utilizing the so-called sudden decay approximation [19, 27]. In a recent paper [25] we considered relaxing the assumptions behind these approximations by allowing for time dependent interactions while evolving the full large-scale perturbation equations. Such an approach can better model the micro physics behind a particular physical framework by allowing one to choose the strength and the time at which the interaction is turned on. In contrast, if the interaction between is modeled with a constant interaction term, the fluid begins to decay (or interact) when its decay width is of the order of the Hubble rate, \( \Gamma \sim H \). Physical scenarios relevant to having time (and space) dependent interactions include e.g. phase transitions, multiple inflation scenarios [33] and scenarios where locally different decay rates of the inflaton are generated by spatially varying reheating temperature and couplings [15, 16, 17].

In the present paper we consider the curvaton scenario with time dependent interactions between curvaton and other fluids. During inflation the curvaton is a light scalar field that does not contribute to the expansion of the universe; after the inflaton field has decayed into relativistic particles the curvaton begins to oscillate and to decay into radiation and matter. The focus of this article is on this situation: we study how a time-dependent interaction affects the evolution of the curvature perturbation. We study the physically allowed parameter space by utilizing information from known cosmological epochs. Moreover, we calculate the amount isocurvature in terms of curvaton

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II. PERTURBATION EQUATIONS

We begin by presenting the equations of motion of the background variables and different density perturbations, where we adapt the notations and conventions of [1, 4] as well as we consider linear scalar perturbations about a spatially flat Friedmann-Robertson-Walker -background in a Newtonian gauge:

\[-ds^2 = g_{\mu\nu}dx^\mu dx^\nu = (1 + 2\phi)dt^2 - a(t)^2(1 - 2\psi)\delta_{ij}dx^i dx^j,\]

where \(a\) is the scale factor. We have used units where \(8\pi G/3 = 1\) and \(c = 1\).

The background evolution is determined by the Einstein’s equations, \(G_{\mu\nu} = 8\pi G T_{\mu\nu}\), and the continuity equations of individual fluids:

\[\dot{\rho}_{(a)} = -3H(1 + \omega_{(a)})\rho_{(a)} + Q_{(a)},\]

where \(\omega_{(a)} = P_{(a)}/\rho_{(a)}\) is the equation of state and \(Q_{(a)}\) describes the energy transfer between different fluids.

Perturbing the covariant continuity equations, one finds the evolution equations of the energy and pressure density perturbations \(\delta\rho_{(a)}\) and \(\delta P_{(a)}\) on large scales [4]:

\[\dot{\delta}\rho_{(a)} + 3H(\delta\rho_{(a)} + \delta P_{(a)}) - 3(\rho_{(a)} + P_{(a)})\dot{\psi} = Q_{(a)}\phi + \delta Q_{(a)}.\]

In addition, we have the \(G_0^0\) component of perturbed Einstein equations

\[3H(\dot{\psi} + H\phi) = -4\pi G\delta\rho,\]

where \(\delta\rho = \sum_{a} \delta\rho_{(a)}\). For perfect fluids, \(\phi = \psi\), and hence given the equations of state, \(\omega_{(a)}\), and the interactions between the fluids, \(Q_{(a)}\), one can evolve the individual fluid perturbations along with the metric perturbation \(\phi\).

Since gauge dependence is an important issue in perturbation analysis we have used gauge invariant perturbations \(\xi_{(a)}\), which equal density perturbations \(\delta\rho_{(a)}\) in uniform curvature gauge \(\psi = 0\), i.e.

\[\xi_{(a)} = \delta\rho_{(a)} + \rho'_{(a)}\psi = -\rho'_{(a)}\zeta_{(a)},\]

\[\zeta = \sum_{a} \xi_{(a)},\]

where \(\zeta_{(a)}\) is the curvature perturbation in uniform density gauge, i.e. \(\delta\rho_{(a)} = 0\) and comma represents derivative with respect to the number of e-folds \(N\) ln \(a\).

The corresponding equations of motion for the gauge invariant quantities are thus

\[\xi'_{(a)} = -3(1 + \omega_{(a)})\xi_{(a)} - 3\delta P_{int(a)} + \frac{\delta Q_{(a)}}{H} + \frac{Q'_{(a)}}{H} - \frac{Q_{(a)}}{H} \left(\frac{\xi}{2\rho_0}\right)\]

\[= -3(1 + \omega_{(a)})\xi_{(a)} - 3\delta P_{int(a)} + \frac{1}{H} \left(\sum_{c} \frac{\partial Q_{(a)}}{\partial \rho_{(c)}}\xi_{(c)} + \frac{\partial Q_{(a)}}{\partial \phi}\phi\right) - \frac{Q_{(a)}}{H} \left(\frac{\xi}{2\rho_0}\right),\]

where \(\delta P_{int(a)} \equiv \delta P_{(a)} - p'_{(a)}\delta\rho_{(a)}/\rho'_{(a)}\) are the internal pressure perturbations. We have also included the possibility of explicit time dependence of the interaction terms. The equation of motion of the curvature perturbations can be derived from eqs. [13] and [21]. The result is

\[\xi'_{(a)} = \frac{3\delta P_{int(a)}}{\rho_{(a)}} - \frac{1}{H\rho'_{(a)}}\delta Q_{int(a)} - \frac{H'\rho_{(a)}}{H\rho_{(a)}}(\zeta - \zeta_{(a)}),\]

where \(\delta Q_{int(a)} \equiv \delta Q_{(a)} - Q'_{(a)}\delta\rho_{(a)}/\rho'_{(a)}\).
The model in hand is a three fluid system with curvaton, radiation and non-relativistic matter fields denoted by subscripts $\sigma$, $\gamma$ and $m$, respectively. Since the curvaton field undergoes coherent oscillations, the curvaton fluid can be safely estimated to behave like non-relativistic matter. Thus we have $\omega_\sigma = 0$, $\omega_\gamma = 1/3$ and $\omega_m = 0$ and the background equations \cite{2} can be written in terms of fractional densities $\Omega_a \equiv \rho_a / \rho$ for which the equations of motion are \cite{3}:

\[ \Omega'_\sigma = \Omega_\sigma \Omega_\gamma + \frac{Q_\sigma}{H \rho}, \]
\[ \Omega'_\gamma = \Omega_\gamma (\Omega_\gamma - 1) + \frac{Q_\gamma}{H \rho}, \]
\[ \Omega'_m = \Omega_m \Omega_\gamma + \frac{Q_m}{H \rho}, \]
\[ \left( \frac{1}{H} \right)' = \left( 1 + \frac{1}{3} \Omega_\gamma \right) \left( \frac{1}{H} \right). \]  

From the definition of $\Omega_a$ it can be easily seen that $\Omega_\sigma + \Omega_\gamma + \Omega_m = 1$, which means that one of the equations of motion is redundant. The interaction terms read as

\[ Q_\sigma = -\Gamma_\gamma f_\gamma(N) \rho_\sigma - \Gamma_m f_m(N) \rho_\sigma, \]
\[ Q_\gamma = \Gamma_\gamma f_\gamma(N) \rho_\sigma, \]
\[ Q_m = \Gamma_m f_m(N) \rho_\sigma, \]  

where $\Gamma(a)$ denotes the strength of the interaction and functions $f(a)(N)$ include the explicit $N$-fold (time) dependence.

From eq. \cite{4} we can finally derive the equations governing the evolution of the perturbations of curvaton-radiation-matter system. The equations for gauge invariant perturbations are

\[ \xi'_\sigma = -3 \xi_\sigma - \frac{\Gamma_\gamma f_\gamma(N) + \Gamma_m f_m(N)}{H} \xi_\sigma - \frac{\Gamma_\gamma f'_\gamma(N) + \Gamma_m f'_m(N)}{H \rho_\sigma} \rho_\sigma \phi - \frac{Q_\sigma}{H} \frac{\xi}{2 \rho}, \]
\[ \xi'_\gamma = -4 \xi_\gamma + \frac{\Gamma_\gamma f_\gamma(N)}{H} \xi_\sigma - \frac{\Gamma_\gamma f'_\gamma(N)}{H} \rho_\sigma \phi - \frac{Q_\gamma}{H} \frac{\xi}{2 \rho}, \]
\[ \xi'_m = -3 \xi_m + \frac{\Gamma_m f_m(N)}{H} \xi_\sigma + \frac{\Gamma_m f'_m(N)}{H} \rho_\sigma \phi - \frac{Q_m}{H} \frac{\xi}{2 \rho}, \]
\[ \phi' = \left( \frac{H'}{H} - 1 \right) \phi - \frac{\xi}{2 \rho}. \]

Correspondingly, the equations for curvature perturbations $\zeta_a$ in the uniform density gauge read as

\[ \zeta'_\sigma = \frac{\Gamma_m f'_m(N) \rho_\sigma (\psi + \zeta_\sigma)}{H \rho_\sigma} + \frac{H'}{H} \left( \Gamma_\gamma + \Gamma_m f_m(N) \right) (\zeta - \zeta_\sigma), \]
\[ \zeta'_\gamma = \frac{\Gamma_\gamma \rho_\sigma (\zeta_\sigma - \zeta_\gamma)}{H \rho_\gamma} - \frac{H'}{H} \frac{\Gamma_\gamma \rho_\sigma}{\rho_\gamma} (\zeta - \zeta_\gamma), \]
\[ \zeta'_m = \frac{\Gamma_m}{H \rho_m} \left[ f_m(N) \rho_\sigma (\zeta_m - \zeta_\sigma) + f'_m(N) \rho_\sigma (\psi + \zeta_m) \right] - \frac{H'}{H} \frac{\Gamma_m f_m(N) \rho_\sigma}{\rho'_m} (\zeta - \zeta_\sigma). \]

The evolution of the system described by eqs \cite{8} and \cite{10} is then ready to be solved numerically once the initial conditions have been set.

Because some of the initial conditions lead to non-physical solutions (e.g. too high reheating temperature or the universe might become matter dominated during nucleosynthesis), the numerical analysis has to be performed carefully. In order to eliminate clearly unphysical scenarios we use physical knowledge from notable epochs of cosmology, namely Big Bang Nucleosynthesis (BBN), radiation-matter equality and decoupling. By fixing the temperatures at those times we are able to determine constraints for the energy distribution between different components of the system. The physical temperature scales have been set to coincide with values given in \cite{32}, e.g. nucleosynthesis came about at temperature at least 0.1 MeV, the matter-radiation equality was reached when $T = 1.0$ eV and decoupling occurred...
when $T = 0.1$ eV. Because $\rho_m, \rho_\sigma \propto a^{-3}$ and $\rho_\gamma \propto a^{-4}$ we can derive a lower limit for the abundances of curvaton, CDM and radiation during nucleosynthesis. Straightforward calculation results in the limit

$$\frac{\Omega_\gamma}{\Omega_\sigma + \Omega_m} \bigg|_{\text{nuc}} \geq 10^5. \tag{12}$$

This limit alone does not, however, guarantee that the curvaton has decayed before the nucleosynthesis. Because even a small contribution of curvaton during nucleosynthesis may eventually begin to dominate the system afterwards and therefore lead to a system with undesirable physical properties, we require that it decays rapidly enough to become subdominant. In practice we compel the beginning of nucleosynthesis to happen after the curvaton starts to effectively decay, i.e.

$$N_{\Omega_{s,\text{max}}} \leq N_{\text{nuc}}. \tag{13}$$

where $N_{\Omega_{s,\text{max}}}$ is the time of maximum curvaton proportion. Setting the temperature at radiation-matter equality allows us to estimate the initial temperature of the system, i.e. reheating temperature by $T_{RH} \approx \left(\frac{\rho_{\gamma, RH}}{\rho_{\gamma, eq}}\right)^{1/4}T_{eq}$. More precisely, the reheating temperatures we quote here are lower bounds since the decay of the curvaton adds to the radiation energy density of the system, but numerical work shows that the effect is small and in practice the initial temperature can well be approximated by the above formula.

A. Non-gaussianity

Single field inflation models usually predict the amount of non-Gaussianity to be too low to be detectable by future CMB-surveys. In contrast multiple scalar fields (e.g. the curvaton scenario) can lead to a clearly observable non-Gaussianity being therefore testable in near future. The non-Gaussianity parameter $f_{NL}$ \cite{7, 27} is defined via gauge invariant Bardeen potential

$$\Phi = \Phi_g + f_{NL} \Phi_g^2, \tag{14}$$

where $\Phi_g$ is the gaussian part of $\Phi$. The Bardeen potential can be expressed at the time of decoupling when the universe is matter dominated as $\Phi = \frac{3}{5}\zeta$. The non-Gaussianity parameter $f_{NL}$ is common to express utilizing an additional parameter defined by

$$r = \frac{\zeta}{\zeta_{\sigma,0}} \bigg|_{\text{dec}}, \tag{15}$$

which tells how effectively the initial curvature perturbation transfers to the total curvature perturbation. Thereby, using the definition of the curvature perturbation $\zeta$, $r$-parameter, the equation of motion of $\rho_\sigma$ and a second order estimate $\delta \rho_\sigma/\rho_\sigma = 2\delta \sigma/\sigma + (\delta \sigma/\sigma)^2$ \cite{27}, the Bardeen potential can be cast in the form

$$\Phi = \frac{2}{5} \frac{r}{1 + \frac{r^2}{3H_0^2}} \left(\frac{\delta \sigma}{\sigma}\right)_0 + \frac{5}{4} \left(\frac{2}{5} + \frac{r}{3H_0} \left(\frac{\delta \sigma}{\sigma}\right)_0\right)^2 \tag{16}$$

and hence $f_{NL} \approx 5/(4r)$. Since during decoupling $\zeta \simeq \zeta_{m,\text{Dec}}$ and in the three-fluid curvaton model $\zeta_m|_{\text{Dec}} = \zeta_{\sigma,0}$ \cite{6}, our value for the transfer parameter is 1 and $f_{NL} = 5/4$. As pointed out in \cite{27} the previous estimate of $f_{NL}$ is valid only when $f_{NL} \gg 1$ since we are using first-order perturbation theory and we are assuming second-order terms $\Phi^{(2)}$ to be at most of order $\Phi_g^2$. Thus the three-fluid model is unable to give any limitations on the non-Gaussianity parameter when the linear theory is applied.

This result differs from the non-gaussianity given in \cite{6} mainly because our $f_{NL}$ is defined at the time of decoupling whereas in \cite{6} the time of nucleosynthesis is used. Our analysis, however, follows the general formalism presented in \cite{31} which allows our $f_{NL}$ to be compared to the observable first order Sachs-Wolfe effect \cite{20}

$$\frac{\Delta T}{T} = \phi + \frac{\delta \rho_\gamma}{4\rho_\gamma}, \tag{17}$$

which is evaluated at the last scattering surface and where the lapse function $\phi = -\Phi$ in the Newtonian gauge.
IV. NUMERICAL RESULTS

We have studied the evolution of equations (8) and (10) in two physically distinct scenarios: A the curvaton decays promptly into radiation and matter i.e. $f_\gamma(N) = 1$, $f_m(N) = 1$ and B curvaton decays first only into radiation component and the matter interaction begins when $N = N^*$ i.e. $f_\gamma(N) = 1$, $f_m(N) = \theta(N - N^*)$. In both scenarios we have systematically scanned the parameter space in order to identify physically acceptable parameter values.

A. Continuous interactions, $f_\gamma(N) = 1$, $f_m(N) = 1$

We have rigorously searched for initial values for which the system is physically motivated i.e. it passes the tests.

FIG. 1: Curvaton-radiation-matter system with different values of $\Gamma_\gamma$ and $\Gamma_m$, when (a) $\Omega_{\sigma0} = 10^{-2}$, (b) $\Omega_{\sigma0} = 10^{-5}$, (c) $\Omega_{\sigma0} = 10^{-7}$ and (d) $\Omega_{\sigma0} = 10^{-10}$. Thick lines represent the initial system temperature in units of GeV.

We have rigorously searched for initial values for which the system is physically motivated i.e. it passes the tests.
mentioned in the previous section. This case has been previously studied in [6]. In figures 1(a)-1(d) these tests can be seen as shades of black and white. A system that passes all of the above mentioned tests is labeled by white color whereas the opposite case is black. Shades of gray indicate that some but not all of the aforementioned physical requirements have been fulfilled. Besides these tests we have also included different contours of the initial system temperature into the figures.

From the figures, we can read out that smaller values of initial curvaton density in general indicate higher initial temperatures, or higher reheat temperatures. The highest reheat temperatures correspond to higher values of \( \Gamma_m \), i.e. larger portion of the curvaton decays into radiation and therefore pushes the time of radiation-matter equality later leading to higher initial temperatures. Since the model applies only when the curvaton field is oscillating, this means that the temperature during reheating has been even higher. However as can be seen from figures, decreasing \( \Omega \) an additional area becomes allowed corresponding to larger values of \( \Gamma_m \). This is a result of the fact that if \( \Gamma_\gamma \) and \( \Gamma_m \) are too small the curvaton does not decay fast enough, leading to possible issues during nucleosynthesis as mentioned above (e.g. for \( \Omega_{\sigma 0} = 10^{-10}, \Gamma_\eta / H_0 = 10^{-20} \) and \( \Gamma_m / H_0 = 10^{-20} \)). By increasing \( \Gamma_m \) the curvaton field decays before nucleosynthesis and therefore leads to a physically sound system (e.g. for \( \Omega_{\sigma 0} = 10^{-10}, \Gamma_\eta / H_0 = 10^{-20} \) and \( \Gamma_m / H_0 = 10^{-6} \)).

The figures indicate that if \( \Gamma_{\gamma,m} \sim H \), a small initial curvaton density is required in order for the system to be physically viable. In addition, we see that in this case the reheat temperature is typically quite low.

### B. Time dependent interactions, \( f_\gamma(N) = 1, f_m(N) = \theta(N - N_*) \)

When the curvaton-matter interaction is turned on at \( N_* \), the system is driven towards an equilibrium of \( \zeta_\gamma \rightarrow -\sigma(N_*) \). This can be seen clearly from the equation of motion of \( \zeta_\gamma \), i.e. eq. (11), because at times close to \( N_* \), terms involving \( f_m'(N) \) dominate the evolution and \( \rho_\gamma < 0 \). Thus the time dependent scenario resembles the situation we previously studied in [26]. As can be seen from figures 2(a)-2(d) different initial values lead to allowed regions very similar to the previous one. Major alteration comes from the change of \( H_0 \rightarrow H(N_* < H_0) \), which shifts different regions upward compared to the time independent scenario. Since the studied cases are initially radiation dominated, the Friedmann equation \( H^2 = \rho_0 \approx \rho_\gamma + \rho_\eta = \rho_{\sigma 0} e^{-3N} + \rho_{\gamma 0} e^{-4N} \approx \rho_{\gamma 0} e^{-4N} \) allows us to estimate \( H(N_*)/H_0 \sim 10^{-7} \) when \( N_* = 8 \). If the initial contribution of curvaton will be larger the system will become curvaton dominated earlier and therefore lead to a smaller value of \( H(N_* \).

Another change can be seen in the initial system temperatures which rise steeply when \( \Gamma_\gamma \) increases. Still, higher values of \( \Gamma_\gamma \) lead to a larger contribution of radiation as in the time independent case. However, since now the matter interaction is turned on later, the curvaton might decay completely to radiation giving very low or no matter contribution at all. This in turn pushes initial temperatures up. Thus again we see that the initial curvaton density needs to be small. But a difference to the continuous-interaction case is that now the reheat temperature can be much higher.

### C. Isocurvature

Isocurvature is defined as the difference between two curvature perturbations \( S_{m \gamma} = 3(\zeta_m - \zeta_\gamma) \). Since the matter curvature perturbation \( \zeta_m \) is here ultimately always driven to the value \( \zeta_{\gamma 0} \) the amount of isocurvature depends only on the final value of \( \zeta_\gamma \). The behavior of the radiation perturbation in the three-fluid model has been studied previously in [3]. They found that if the curvaton begins to dominate before the decay epoch almost all of the radiation originates from the curvaton fluid and therefore gives \( \zeta_\gamma \approx \zeta_{\gamma 0} \).

We have plotted the amount of isocurvature in the three-fluid model in figures 3(a) and 3(b). Our results clearly agree with the reasoning above. If \( \Gamma_m \) or \( \Gamma_\gamma \) is large enough, the curvaton fluid will decay before it begins to dominate the system and hence lead to a non-vanishing final isocurvature. For a smaller initial curvaton, it takes even longer for the curvaton fluid to dominate and therefore the system is adiabatic only for very small values of \( \Gamma_m \) and \( \Gamma_\gamma \). These in turn lead to high reheating temperatures, e.g. for initial values \( \Omega_{\sigma 0} = 10^{-10}, \Gamma_\eta / H_0 = 10^{-26} \) and \( \Gamma_m / H_0 = 10^{-34} \) a numerical evaluation gives a reheating temperature \( 10^{12} \text{ GeV} \) and \( S_{m \gamma}/\zeta_{\text{dec}} = 0.0006 \). The time dependent interaction gives very similar results once the scaling \( \Gamma_m / H_0 \rightarrow \Gamma_m / H_* \) is taken into account.
FIG. 2: Curvaton-radiation-matter system where matter interaction is turned on at $N = 8$ with different values of $\Gamma_\gamma$ and $\Gamma_m$, when (a) $\Omega_{\sigma 0} = 10^{-2}$, (b) $\Omega_{\sigma 0} = 10^{-5}$, (c) $\Omega_{\sigma 0} = 10^{-7}$ and (d) $\Omega_{\sigma 0} = 10^{-10}$. Thick lines represent the initial system temperature in units of GeV.

V. DISCUSSION AND CONCLUSIONS

The curvaton model has gained a lot of attention in the recent years mainly because it can make the inflation potential look more natural [19]. Since the first model where the curvaton decayed only into radiation [19, 21], a number of other possibilities have been explored including a curvaton web model [22], the possibility of multiple curvaton fields [23] and different particle models such as axions [24, 25] just to name a few.

In the present paper we have studied a three-fluid model in which the curvaton decays into both radiation and cold dark matter. This has been studied previously also in [6, 27, 28]. We have assumed that the initial system has no matter content and it is dominated by the radiation which originates from the decay of the inflaton field. Because all of the matter content of the universe comes from the curvaton field we are able to estimate the reheating temperature.
This can be additionally used to constraint the parameter space of the model.

We have systematically scanned the parameter space and identified the regions where the model is physically acceptable, i.e., when evolution during and after nucleosynthesis is standard while requiring that the reheating temperature is not unreasonably high. We have identified these regions both when the decay rates are fixed and when the curvaton starts to decay later. These allowed regions are alike once the rescaling $\Gamma / H_0 \rightarrow \Gamma / H(N_\ast)$ is taken into account and the real difference appears in the initial system temperature.

We find that if the decay rates are comparable to the Hubble rate, a small initial curvaton density is required. Otherwise one needs to fine-tune the decay rates to be much smaller than $H$ at the time of decay. In the continuous interaction case, requiring $\Gamma_i \sim H$ leads to a low reheat temperature, but this can be avoided when the matter interaction is delayed.

If the initial curvaton density is large, the final state is naturally adiabatic assuming that the system is otherwise physically acceptable. Note however, that this in turn requires fine-tuning in the decay rates. If $\Gamma_i \sim H$, we find that the final state generally contains a large isocurvature component.

We have also studied non-gaussianity in the framework of the three-fluid models. We find that in the region where the first-order perturbation theory can be applied, the three-fluid model gives no limits on the $f_{NL}$ parameter. This is the result of a conserved curvature perturbation $\zeta_c$, which carries the initial curvaton perturbation $\zeta_\sigma$ into the matter perturbation $\zeta_m$ [6]. Our results differ from the previous results [6, 27] mainly because our $f_{NL}$ is evaluated at the time of last scattering and not at nucleosynthesis. This allows the non-gaussianity to be more easily compared to the observational Sachs-Wolfe effect and we do not have to use a radiation transfer function. In order to calculate the observational non-gaussianity, second-order perturbation theory needs to be applied [31]. This is however beyond the scope of this article and will be the focus of a follow-up paper.
This project has been partly funded by the Academy of Finland project no. 811953. TM and JS are supported by the Academy of Finland.

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