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Isocurvature cosmological perturbations and the CMB

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Abstract

We present a short review on cosmological isocurvature perturbations, which some early universe models predict in addition to the standard adiabatic perturbation. After defining the isocurvature perturbations, we analyse their evolution within the framework of the standard perturbation theory. We then explain how these perturbations can be constrained by the CMB data. We finally discuss some mechanisms that generate isocurvature perturbations. ***To cite this article: D. Langlois, C. R. Physique 4 (2003).***

Résumé

Perturbations cosmologiques isocourbures et le CMB. Nous donnons une introduction aux perturbations cosmologiques isocourbures, qui sont prédites par certains modèles de cosmologie primordiale, en plus des perturbations adiabatiques habituelles. Après avoir défini la notion de perturbations isocourbures, nous analysons leur évolution dans le cadre de la théorie standard des perturbations cosmologiques. Nous expliquons ensuite comment ces perturbations peuvent être contraintes par les observations du CMB. Finalement, nous présentons quelques mécanismes qui engendrent des perturbations isocourbures. ***Pour citer cet article : D. Langlois, C. R. Physique 4 (2003).***

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Mots-clés : Anisotropies du fond de rayonnement cosmique ; Cosmologie ; Univers primordial

1. Introduction

It is usually not emphasized enough that the CMB anisotropies are sensitive not only to the homogeneous matter content of the Universe, but also to the initial conditions for cosmological perturbations. CMB anisotropies thus represent an important observational test for any early universe mechanism explaining the origin of cosmological fluctuations. With the increasing precision of the available CMB data, it becomes possible to refine the constraints on the 'primordial fluctuations', and to investigate whether they deviate from the standard adiabatic assumption.

To be precise, 'initial' or 'primordial' perturbations are defined deep in the radiation era but at temperatures low enough, i.e., typically after nucleosynthesis, so that the main cosmological components reduce to the usual photons, baryons, neutrinos and cold dark matter (CDM). Moreover, only modes with wavelength larger than the Hubble radius in this primordial era will be relevant today for CMB anisotropies.

The above various cosmological species can be characterized by their number density, n_X, and their energy density ρ_X. Linear perturbations of these quantities are defined as

$$\delta n_X = n_X - \bar{n}_X, \qquad \delta\rho_X = \rho_X - \bar{\rho}_X, \tag{1}$$

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where the bar corresponds to the homogeneous (unperturbed) quantity.

The *adiabatic mode* is defined as a perturbation affecting all the cosmological species such that the relative ratios in the number densities remain unperturbed, i.e., such that

$$\delta(n_X/n_Y) = 0. \tag{2}$$

It is associated with a curvature perturbation, via Einstein's equations, since there is a global perturbation of the matter content. This is why the adiabatic perturbation is also called *curvature* perturbation. In terms of the energy density contrasts, defined by

$$\delta_X \equiv \frac{\delta\rho_X}{\rho_X}, \tag{3}$$

the adiabatic perturbation is characterized by the relations

$$\frac{1}{4}\delta_\gamma = \frac{1}{4}\delta_\nu = \frac{1}{3}\delta_b = \frac{1}{3}\delta_c. \tag{4}$$

They follow directly from the prescription (2), each coefficient depending on the equation of state of the particuler species.

Since there are several cosmological species, it is also possible to perturb the matter components without perturbing the geometry. This corresponds to *isocurvature* perturbations, characterized by variations in the particle number ratios but with vanishing curvature perturbation. The variation in the relative particle number densities between two species can be quantified by the so-called *entropy perturbation*

$$S_{A,B} \equiv \frac{\delta n_A}{n_A} - \frac{\delta n_B}{n_B}. \tag{5}$$

When the equation of state for a given species is such that $w \equiv p/\rho = \text{Const}$, then one can re-express the entropy perturbation in terms of the density contrast, in the form

$$S_{A,B} \equiv \frac{\delta_A}{1+w_A} - \frac{\delta_B}{1+w_B}. \tag{6}$$

It is convenient to choose a species of reference, for instance the photons, and to define the entropy perturbations of the other species relative to it:

$$S_b \equiv \delta_b - \frac{3}{4}\delta_\gamma, \tag{7}$$

$$S_c \equiv \delta_c - \frac{3}{4}\delta_\gamma, \tag{8}$$

$$S_\nu \equiv \frac{3}{4}\delta_\nu - \frac{3}{4}\delta_\gamma, \tag{9}$$

thus define respectively the *baryon isocurvature mode*, the *CDM isocurvature mode*, and the *neutrino isocurvature mode*. In terms of the entropy perturbations, the adiabatic mode is obviously characterized by $S_b = S_c = S_\nu = 0$.

In summary, we can decompose a general perturbation, described by four density contrasts, into one adiabatic mode and three isocurvature modes. In fact, the problem is slightly more complicated because the evolution of cosmological perturbations is governed by second order differential equations and a perturbed (perfect) fluid is described locally by its density contrast and by its velocity field. The 'primordial' perturbations are constrained by the requirement that the perturbations do not diverge when going backwards in time deep in the radiation era. With this prescription, there remains one arbitrary relative velocity between the species, which gives an additional mode, usually named the *neutrino isocurvature velocity* perturbation.

At this stage, it is worth warning the non-expert reader about the somewhat loose terminology used in the literature. One usually refers to an isocurvature mode with the meaning that this mode was 'initially' an isocurvature mode, i.e., deep in the radiation era. But this 'primordial' isocurvature mode, when considered at later times, for instance at last scattering or today, can have an adiabatic component, because the decomposition between adiabatic and isocurvature is not time-invariant. A 'primordial' pure isocurvature perturbation can generate later an adiabatic contribution if the energy densities of the various species evolve differently so that the balance that ensured an unpertubed total energy density is lost.

The CMB is a powerful way to study isocurvature perturbations because (primordial) adiabatic and isocurvature perturbations produce very distinctive features on the CMB anisotropies. Whereas an adiabatic initial perturbation generates a cosine oscillatory mode in the photon–baryon fluid, leading to an acoustic peak at $\ell \simeq 220$ (for a flat universe), a pure isocurvature initial perturbation generates a sine oscillatory mode resulting in a first peak at $\ell \simeq 330$.

The unambiguous observation of the first peak at $\ell \simeq 220$ has eliminated the possibility of a dominant isocurvature perturbation. The recent observation by WMAP of the CMB polarization has also confirmed that the initial perturbation is mainly an adiabatic mode. However, this does not exclude the presence of a subdominant isocurvature contribution, which could be detected in future high-precision experiments such as Planck.

On the theoretical side, there has been recently a growing interest in *correlated* mixtures of adiabatic and isocurvature perturbations, which can be generated for instance in multiple-field inflation. The detection of a (correlated or uncorrelated) isocurvature mode, if any, would be an invaluable clue to understand the physics of the early universe.

2. Theory of cosmological perturbations

2.1. Evolution equations

We now briefly recall the basic equations derived in the theory of cosmological perturbations (see, e.g., [1]), which are useful to define precisely the primordial isocurvature perturbations and to follow their evolution during the history of the universe.

Let us start with the scalar perturbations of the spacetime geometry, which can be described by two scalar potentials, Φ and Ψ, so that the perturbed (flat) Friedmann–Lemaître–Robertson–Walker (FLRW) metric reads

$$\mathrm{d}s^2 = -a^2(\eta)(1+2\Phi)\,\mathrm{d}\eta^2 + a^2(\eta)(1-2\Psi)\delta_{ij}\,\mathrm{d}x^i\,\mathrm{d}x^j. \tag{10}$$

This description corresponds to the so-called longitudinal gauge. We have introduced the conformal time η rather than the cosmic time t, the two being related by $\mathrm{d}t = a\,\mathrm{d}\eta$.

Instead of the energy density contrasts δ_A defined in the longitudinal gauge (10), it is convenient to introduce the slightly redefined energy density contrasts (defined in the flat-slicing gauge):

$$\Delta_\gamma = \delta_\gamma - 4\Psi, \quad \Delta_\nu = \delta_\nu - 4\Psi, \quad \Delta_b = \delta_b - 3\Psi, \quad \Delta_c = \delta_c - 3\Psi. \tag{11}$$

Let us now present the system of equations governing the evolution of the matter perturbations, which involve the density contrasts of the four cosmological species and their peculiar velocity potentials, denoted by V_A. For convenience, we work directly in the Fourier space rather than in ordinary space, according to the definition

$$f_{\mathbf{k}} = \int \frac{\mathrm{d}^3\mathbf{x}}{(2\pi)^{3/2}}\,\mathrm{e}^{-i\mathbf{k}\cdot\mathbf{x}} f(\mathbf{x}). \tag{12}$$

To make the notation lighter, we will drop, in the following, the subscript $\mathbf{k}$ for the various perturbations.

We first have four equations of conservation, one for each species, that can be written (in Fourier space):

$$\begin{aligned}
\dot\Delta_\nu &= \frac{4}{3}kV_\nu,\\
\dot\Delta_c &= kV_c,\\
\dot\Delta_\gamma &= \frac{4}{3}kV_{b\gamma},\\
\dot\Delta_b &= kV_{b\gamma},
\end{aligned} \tag{13}$$

where a dot denotes a derivative with respect to the conformal time. Since the baryon and photon fluids are strongly coupled via Thomson scattering before the decoupling, one can take a velocity, $V_{b\gamma}$, common to the baryon and photon fluids. After the decoupling, the fluid description no longer applies and one must resort to the Boltzman equation to solve for the evolution of the photon gas.

We then have three Euler equations, two for the independent fluids of CDM and neutrinos, and one for the baryon–photon fluid:

$$\begin{aligned}
\dot V_\nu &= -k\left[\frac{\Delta_\nu}{4} + \Psi + \Phi - \sigma_\nu\right],\\
\dot V_c &= -\mathcal{H}V_c - k\Phi,\\
\dot V_{b\gamma} &= -\frac{3\Omega_b}{4\Omega_\gamma + 3\Omega_b}\mathcal{H}V_{b\gamma} - k\frac{4\Omega_\gamma}{4\Omega_\gamma + 3\Omega_b}\left[\frac{\Delta_\gamma}{4} + \Psi\right] - k\Phi,
\end{aligned} \tag{14}$$

where $\mathcal{H}$ is the comoving Hubble parameter defined by $\mathcal{H} \equiv a'/a$. In the last equation, the coefficients Ω_A are time dependent since the ratio of the energy density of a given species with respect to the critical energy density, will change with time.

To close the above system of equations, one needs the Einstein equations, which express the metric perturbations in terms of the matter perturbations. Only two components of the Einstein equations are useful, the others being redundant, and they are the Poisson equation

$$-\left[\frac{k^2}{\mathcal{H}^2}+\frac{9}{2}(1+w)\right]\Psi=\frac{3}{2}\sum_X \Omega_X\left[\Delta_X-\frac{3\mathcal{H}}{k}(1+w_X)V_X\right], \tag{15}$$

and the anisotropic stress equation,

$$\frac{k^2}{\mathcal{H}^2}(\Psi-\Phi)=6\Omega_\nu\sigma_\nu, \tag{16}$$

where σ_ν represents the anisotropic stress due to the neutrinos (which thus require a decription beyond the perfect fluid approximation). To get the evolution of σ_ν, one must use a higher moment of the Boltzmann equation, which gives

$$\dot{\sigma}_\nu=-\frac{4}{15}kV_\nu, \tag{17}$$

where other terms on the right-hand side have been neglected.

The above equations govern (at least in the fluid approximation) the evolutions of the perturbations. Before using them to define the primordial perturbations and study their subsequent evolution, it is instructive to discuss their statistical properties.

2.2. Power spectra and correlations

In cosmology, perturbations are treated as homogeneous and isotropic random fields. Primordial perturbations are usually assumed to be Gaussian, in which case their statistical properties can be summarized simply in terms of their power spectrum, defined for a quantity f by

$$\langle f_{\mathbf{k}} f^*_{\mathbf{k}'}\rangle = 2\pi^2 k^{-3}\mathcal{P}_f(k)\delta(\mathbf{k}-\mathbf{k}'). \tag{18}$$

When primordial perturbations are described by *several quantities*, such as is the case with mixtures of adiabatic and isocurvature perturbations, one must also define, for any pair of random fields f and g, a *cross-correlation spectrum* $\mathcal{P}_{f,g}(k)$ by the following expression (see [2,3])

$$\Re\langle f_{\mathbf{k}} g^*_{\mathbf{k}'}\rangle = 2\pi^2 k^{-3}\mathcal{P}_{f,g}(k)\delta(\mathbf{k}-\mathbf{k}'). \tag{19}$$

By renormalizing the cross-correlation spectrum, one may define a *correlation coefficient*, comprised between -1 and 1:

$$\cos\Delta(k)=\frac{\mathcal{C}_{f,g}(k)}{\sqrt{\mathcal{P}_f(k)}\sqrt{\mathcal{P}_g(k)}}. \tag{20}$$

The two extreme values $\cos\Delta=+1$ and $\cos\Delta=-1$ correspond respectively to full correlation and full anti-correlation.

Before the work [2], only independent mixtures of adiabatic and isocurvature modes, i.e., with vanishing correlation, were considered in the literature. This statistical independence means that the quantities Φ and S can be expressed as

$$\Phi=\mathcal{P}_\Phi^{1/2}e_1, \qquad S=\mathcal{P}_S^{1/2}e_2, \tag{21}$$

where e_1 and e_2 are *independent* normalized centered Gaussian random fields (i.e., such that $\langle e_i(\mathbf{k})\rangle=0$, $\langle e_i(\mathbf{k})e_j^*(\mathbf{k}')\rangle = \delta_{ij}\delta(\mathbf{k}-\mathbf{k}')$, for $i,j=1,2$), and where the subscripts $\mathbf{k}$ are implicit. With the assumption (21), one obtains immediately a vanishing correlation.

Clearly, if one takes into account all five possible initial modes, as mentioned in the introduction, the most general (Gaussian) perturbation random field must be described by a 5×5 symmetric correlation matrix, as discussed in [4].

2.3. Long wavelength perturbations

As mentioned earlier, primordial perturbations are defined deep in the radiation era on super-Hubble wavelengths, i.e., such that $k\ll aH$. In order to characterize all primordial modes, one can expand the perturbative quantities in terms of the small parameter $k\eta$, so that

$$X=X^{(0)}+X^{(1)}k\eta+X^{(2)}(k\eta)^2+\cdots. \tag{22}$$

One then substitutes these expansions in the equations governing the evolution of the perturbations. This allows us to express all perturbations in terms of the primordial curvature perturbation $\hat{\Phi}$ and the three primordial entropy perturbations, $\widehat{S}_b$, $\widehat{S}_c$ and $\widehat{S}_\nu$ (ignoring the neutrino isocurvature velocity mode).

During most of their history, the relevant cosmological modes remain outside the Hubble radius ($k \ll aH$). In order to study easily the super-Hubble evolution of adiabatic and isocurvature perturbations, it is convenient to introduce the quantity

$$\zeta = -\Psi - H\frac{\delta\rho}{\dot{\rho}}, \tag{23}$$

defining the curvature perturbation on hypersurfaces of uniform (total) energy density. A similar quantity can be defined for any individual fluid,

$$\zeta_X = -\Psi - H\frac{\delta\rho_X}{\dot{\rho}_X}. \tag{24}$$

Using the evolution equations written above, one can show that the evolution of ζ on super-Hubble scales is given by

$$\dot{\zeta} = -\frac{H}{\rho + P}\left(\delta P - c_s^2\delta\rho\right), \tag{25}$$

where the expression between parenthesis, involving the *total fluid*, can be interpreted as a non-adiabatic pressure. The same expression applies for each of the individual fluids as long as they are non-interacting. For each such individual fluid, ζ_X is conserved on super-Hubble scales, since the equation of state is barotropic. Moreover, since

$$S_{A,B} = 3(\zeta_A - \zeta_B) \tag{26}$$

one finds that the isocurvature perturbation between two non-interacting fluids is conserved on super-Hubble scales.

As for the global ζ, it is not conserved if there are isocurvature perturbations, which make the non adiabatic pressure nonzero. It is possible to compute explicitly the expression of the gravitational potential perturbation during the *matter dominated era* in terms of the primordial perturbations [3]. For an illustrative purpose, we will discuss here only the case of CDM isocurvature perturbations, which is the case the most studied in the literature. In the matter era, the gravitational potential perturbation is given by

$$\Phi_{\text{matter}} = \frac{3}{10}\left(3 + \frac{4}{5}\Omega_\nu^{\text{RD}}\right)\hat{\Phi} - \frac{1}{5}\Omega_c^{\text{MD}} S_c \simeq \hat{\Phi} - \frac{1}{5}S_c, \tag{27}$$

where Ω_ν^{RD} is the energy density fraction of neutrinos in the radiation era and Ω_c^{MD} is the energy density fraction of CDM in the matter era.

Using these results, one can express the CMB anisotropies on large scales (i.e., on scales that are super-Hubble at the time of last scattering) explicitly in terms of the primordial perturbations. One finds

$$\left(\frac{\Delta T}{T}\right) = \frac{1}{10}\left(3 + \frac{4}{5}\Omega_\nu^{\text{RD}}\right)\hat{\Phi} - \frac{2}{5}\Omega_c^{\text{MD}} S_c \simeq \frac{1}{3}\hat{\Phi} - \frac{2}{5}S_c. \tag{28}$$

By comparing the above expressions (27) and (28), one recovers the well-known results that $\Delta T/T \simeq \Phi/3$ for pure adiabatic initial conditions and $\Delta T/T \simeq 2\Phi$ for pure isocurvature perturbations.

3. CMB constraints on isocurvature modes

The first bounds on isocurvature perturbations assumed uncorrelated modes [6]. Recently, however, several works have analysed the CMB data looking for correlated adiabatic-isocurvature perturbations. Considering only the adiabatic mode and *one* isocurvature mode, the following spectra have been assumed:

$$\mathcal{P}_{\mathcal{R}}(k) = \mathcal{A}^2\left(\frac{k}{k_0}\right)^{n_{\text{ad}}-1},$$

$$\mathcal{P}_{\mathcal{S}}(k) = \mathcal{B}^2\left(\frac{k}{k_0}\right)^{n_{\text{iso}}-1}, \tag{29}$$

$$\mathcal{P}_{\mathcal{RS}}(k) = \mathcal{A}\mathcal{B}\cos\Delta\left(\frac{k}{k_0}\right)^{(n_{\text{ad}}+n_{\text{iso}})/2-1}, \tag{30}$$

where k_0 is an arbitrary pivot scale and we have introduced, instead of the gravitational potential perturbation Φ, the curvature perturbation $\mathcal{R}$, often used in the literature, related to Φ during the radiation-dominated era by

$$\mathcal{R} = \frac{3}{2}\Phi. \tag{31}$$

Ignoring the overall amplitude of the fluctuations, the above parametrization means that the initial perturbations are characterized by four parameters: the adiabatic index $n_{\rm ad}$, the isocurvature index $n_{\rm iso}$, the isocurvature/adiabatic amplitude ratio $\mathcal{B}/\mathcal{A}$ and the correlation parameter $\cos\Delta$. This represents three more parameters than the usual description, which includes only the adiabatic index $n_{\rm ad}$.

Allowing for isocurvature perturbations in addition to adiabatic perturbations introduces additional degeneracies in the interpretation of the CMB data. In particular, analysing the CMB data with the prior assumption that primordial perturbations are purely adiabatic could lead to an incorrect determination of the cosmological parameters if the real universe contains an isocurvature contribution. To lift these degeneracies, the measurement of polarisation is crucial, as it was shown quantitatively in [5].

Several very recent papers [7–9] have considered the isocurvature modes in the light of the WMAP data, improving similar previous studies [10,11]. The present situation is that there is no significant improvement of the fit to the data, from which one can conclude that the present data do not require an isocurvature component as parametrized above. In a different and more positive perspective, this means that these analyses provides us with constraints on the maximal contribution of an isocurvature component to the CMB signal, which can be turned into constraints on the early universe models. Note that the constraints on the CDM isocurvature mode are much more stringent than those on the other types of isocurvature modes.

4. Generation of isocurvature perturbations

Isocurvature perturbations have been invoked in many cosmological scenarios. Pure isocurvature initial conditions have even been proposed as an alternative to the standard adiabatic initial conditions, until the observation of the first acoustic peak rules out this idea. Isocurvature perturbations can also play an important role in models inspired by particle physics, such as cosmological axions [12] and Affleck–Dine baryogenesis [13].

A point worth noticing, however, is that, in contrast with the adiabatic mode which can always be seen as a perturbation of the geometry, the isocurvature modes can be erased during the thermal history of the universe. Indeed, if the microphysical processes are such that all species are locally in a thermal equilibrium that depends only on the local temperature T, then all the number densities depend only on T, which excludes any isocurvature perturbation. Therefore, the existence and the evolution of the isocurvature modes are very sensitive to the microphysical processes that took place in the early universe.

Below, we discuss two types of mechanisms than can produce isocurvature perturbations and which have been explored very actively during the last few years.

4.1. Multiple-field inflation

In its simplest version, inflation is driven by a single scalar field, which generates an approximately scale-invariant Gaussian curvature perturbations [14]. This single field then decays into ordinary matter and automatically gives primordial adiabatic perturbations. However, current models describing high energy physics usually contain many scalar fields, which suggests the possibility that several scalar fields play a significant role during inflation.

Multiple-field inflation can generate, in addition to the adiabatic component, an isocurvature contribution [15]. With appropriate conditions, multiple-inflation can even generate *correlated* adiabatic and isocurvature perturbations [2,16,17]. In the case of two uncoupled free massive scalar fields, one can show that the 'primordial' curvature and entropy perturbations, i.e., deep in the radiation era, can be expressed as linear combinations of the scalar field fluctuations at Hubble radius crossing, i.e.,

$$\hat{\Phi} = A_1\delta\phi_1 + A_2\delta\phi_2, \qquad \widehat{S} = B_1\delta\phi_1 + B_2\delta\phi_2. \tag{32}$$

If the weights are such that, say, $A_2 \ll A_1$ and $B_1 \ll B_2$, then one obtains uncorrelated adiabatic and isocurvature perturbations. This is the case for instance if inflation is driven by ϕ_1 while ϕ_2 is a very light scalar field and remains passive during the inflationary phase relevant for current cosmological scales. Then, essentially, the fluctuations of ϕ_1 generate the adiabatic component while the fluctuations of ϕ_2 give the isocurvature component. However, if our cosmological 'window' corresponds to a phase of double inflation when both scalar fields have a dynamical role, then the weights in (32) can be of the same order of magnitude, which leads to strong correlations between the adiabatic and isocurvature perturbations.

It can also be convenient to define adiabatic and isocurvature modes during inflation [16]. These *instantaneous modes* can be defined in terms of the perturbations of the two scalar fields. The adiabatic and isocurvature perturbations can be interpreted as the projections respectively along and orthogonal to the background trajectory. One must not confuse these modes with the 'primordial' adiabatic and isocurvature modes, defined deep in the radiation era. Formally, they can be related by a transformation of the form [11]

$$\begin{pmatrix} \widehat{\mathcal{R}}_{\rm rad} \\ \widehat{\mathcal{S}}_{\rm rad} \end{pmatrix} = \begin{pmatrix} 1 & T_{\mathcal{RS}} \\ 0 & T_{\mathcal{SS}} \end{pmatrix} \begin{pmatrix} \mathcal{R}_* \\ \mathcal{S}_* \end{pmatrix}, \tag{33}$$

where the subscript $*$ denotes horizon crossing during inflation. $T_{\mathcal{RR}} = 1$ since the curvature perturbation is conserved for purely adiabatic perturbations whereas $T_{\mathcal{SR}} = 0$ because adiabatic perturbations cannot source entropy perturbations on large scales. The other terms are model dependent. A potentially observable effect of some multi-field inflation models is the generation of a significant non-Gaussianity in the primordial perturbations, as recently discussed in [18], in contrast with single field inflation [19].

4.2. Curvaton

Another scenario related to isocurvature perturbations has been recently explored: the curvaton scenario [20]. It is based on the fact that a pure isocurvature primordial perturbation can generate a curvature perturbation at late times. The curvaton σ is a scalar field that is very light during inflation and acquires Gaussian fluctuations. These fluctuations will give isocurvature fluctuations. When the Hubble parameter decreases below the curvaton mass, the curvaton oscillates and behaves as pressureless matter. If it decays sufficiently late, it ends up dominating the matter content of the Universe. This is associated with a transformation of the isocurvature perturbations into a curvature perturbation. One gets

$$\widehat{\mathcal{R}} = -\frac{r}{4+3r}\delta_\sigma, \tag{34}$$

where r is the ratio $\rho_\sigma/\rho_{\mathrm{rad}}$ when the curvaton decays. In the simplest version, the curvaton perturbations are totally converted into adiabatic perturbations. One can, however, envisage scenarios where the scenario will also generate an isocurvature component in addition to the curvature one [21]. If this is the case the adiabatic and isocurvature perturbations are fully correlated (or anti-correlated) since they are produced by the same field. The curvaton scenario can also produce important non-Gaussianities in the primordial perturbations.

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