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Primordial Perturbations, Early Evolution and Dark Matter in the Universe*

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Abstract

The problem of primordial cosmological perturbations which gave birth to the large scale structure of the universe is analysed from the time of their origin up to the nonlinear formation of the structure. We show how classical gravity accounts for the generation of perturbations near the beginning of the cosmological expansion. Dark matter which governs the further development of perturbations in the early universe includes both heavy and light weakly interacting particles. Transfer functions of the adiabatic and isothermal modes of perturbations are investigated and some astrophysical applications are given.

1. Introduction

Current cosmology suggests that all the observable structure of our universe on large scales (galaxies, clusters and superclusters) has stemmed from the small primordial perturbations of matter density and gravitational field which had to be present very early in the Friedmann universe. For galaxies to form nowadays we need the primordial perturbation amplitude near the Big Bang to be of the order of $\sim 10^{-4}$. On the one hand, it is small for the linear perturbation theory to be used from the very beginning but, on the other hand, it is many orders of magnitude greater than the 'natural' inhomogeneities (quantum, statistical, thermal etc.) to be expected in the homogeneous universe.

So, at least two problems must be considered before one could simulate galaxy formation at the nonlinear late stages (for redshifts $Z \lesssim 10$): production of the cosmological primordial perturbations in the very early universe, their original spectrum, distribution and other properties; the evolution of the created perturbations in the early universe (until they became large), which has modified and changed their original characteristics.

Further on, I try to outline the principal ideas which modern cosmology proposes to solve the first problem and review some results for the transfer functions which bear all the information about the changes of primordial spectra before perturbations enter the nonlinear evolution. Some astrophysical applications and isocurvature models are also discussed.

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2. Difficulties of the Classical Hot Universe and the Idea of 'Parametric' Amplification

According to the Lifshitz (1946) theory of small perturbations there are three types of perturbation to the homogeneous isotropic universe:

- ⋆ the density (potential) perturbations,
- ⋆ the vortex perturbations,
- ⋆ gravitational waves.

We are interested now in the first type—the potential or scalar perturbations—because we believe that they are responsible for galaxy formation. We also assume that the universe is spatially flat on scales encompassing large structures ($l < 10^3$ Mpc):

$$ds^2 = a^2(d\eta^2 - dx^2 - dy^2 - dz^2),$$

where $a = a(\eta)$ is a scale factor, and the universal time is $t = \int a\, d\eta$ ($c = 8\pi G = \hbar = 1$).

First, let us show that in the hot universe, when the total matter is dominated by relativistic particles and $a \sim \eta$, the production of perturbations is impossible. Indeed, it is well known that in this case the Fourier amplitude of an arbitrary perturbation is easily separated into two modes—the so-called 'growing' and 'decaying' modes. It is convenient to present the exact solution in terms of the gauge invariant scalar q which fully governs the behaviour of the potential perturbations (Lukash 1980):

$$q = C_1 \frac{\sin\kappa}{\kappa} + C_2 \frac{\cos\kappa}{\kappa}, \tag{1}$$

where the functions of wavenumber $C_{1,2} = C_{1,2}(k)$ are the 'growing' and 'decaying' mode amplitudes respectively; $\kappa = k\eta/\sqrt{3} \sim t/\lambda$, and $\lambda = 2\pi a/k$ is the physical wavelength.

All the other physical quantities are easily related to the scalar q. In the synchronous gauge the density and the gravitational potential perturbations can be written (let us recall that the equation of matter state is $p = \epsilon/3$):

$$\frac{\delta\epsilon}{\epsilon} = C_1\left\{-\cos\kappa + 2\left(\frac{\sin\kappa}{\kappa} + \frac{\cos\kappa - 1}{\kappa^2}\right)\right\}$$
$$+ C_2\left\{\sin\kappa + 2\left(\frac{\cos\kappa}{\kappa} - \frac{\sin\kappa}{\kappa^2}\right)\right\},$$
$$h = C_1\left(\frac{1-\cos\kappa}{\kappa^2}\right) + C_2 \frac{\sin\kappa}{\kappa^2}. \tag{2}$$

The most important feature of these perturbations is the following. Neither 'growing' mode nor 'decaying' mode increases catastrophically: both of them are described by sin and cos terms, so if C_1 and C_2 are less than unity both these modes become sound waves with constant (independent of time) amplitudes C_1 and C_2. This result means that the hot universe is absolutely stable against gravitational instability. If the initial perturbations were small they will be small forever (Bisnovatyi-Kogan *et al.* 1980; Lukash and Novikov 1986).

These sound waves with small amplitudes existed in the hot universe until a few hundred years after the beginning of the expansion. At this epoch the universe became cold so that the equation of state $p = \epsilon/3$ is no longer valid. Now the clumps of the medium in the sound waves begin to grow because of the real gravitational instability (at least some components of the medium undergo this instability) and this process develops causing the fragmentation of the medium into separate bodies.

We shall not discuss these late processes of galaxy formation at this stage. The point is that for the formation of galaxies we need a definite amplitude of sound waves of $\sim 10^{-4}$ in the linear scale which encompasses a large-enough number of baryons for galaxy formation. So, C_1 or/and C_2 must be on this scale of the order of 10^{-4}. This is a very stringent demand on the initial perturbations. Indeed, when t is small, $\kappa \ll 1$, we have

$$h_1 \approx C_1/2 \ll 1, \qquad h_2 \approx C_2/\kappa \ll 1 . \tag{3}$$

From these expressions we can see that C_2 must be extremely small and could not be of the order of 10^{-4}. So, we need two conditions to be met:

$$\text{(i) } C_2 \ll C_1; \quad \text{(ii) } C_1 \approx 10^{-4} . \tag{4}$$

Both of them look very strange, and indeed:

(I) In any reasonable assumption about the initial seed perturbations (i.e. for random phase fluctuations) C_1 must be equal to C_2 and (i) could not be correct.

(II) For any 'natural' fluctuations, C_1 and C_2 are dozens of orders of magnitude less than 10^{-4}.

The second conclusion can be clarified by the following example. Let us suppose that the time of origin of the fluctuations is the Planck time $t_0 = t_{\rm Pl}$, and let us denote $k = 1$ for $\lambda_{\rm Pl}$. On the scale of galaxies we have $k_{\rm gal} \approx 10^{-26}$. Now let us suppose (as an example) that the spectrum of the $\delta\epsilon/\epsilon$ fluctuation at that moment had a thermal shape with maximum at $\lambda = \lambda_{\rm Pl}$; then the amplitude of the perturbation would be proportional to $k^{3/2}$, and on the scale of galaxies the amplitude is $k_{\rm gal} \approx 10^{-40}$. Thus C_1 has to be $\sim 10^{-40}$ and so is 35 orders of magnitude less than we need.

Our conclusions are the following: classical cosmology of the hot universe has great difficulty in explaining the primordial seed perturbations of matter. Fluctuations in the classical hot universe require:

(i) $C_1 \gg C_2$—not 'natural',

(ii) $C_1 \approx 10^{-4}$—much greater than the 'natural' value $C_{\rm gal} \approx 10^{-40}$.

So, we come to the conclusion that the way to overcome these difficulties is to reject the 'hot' equation of state $p = \epsilon/3$, i.e. to break the linear law of expansion $a \sim \eta$ at some very early stage. This is how we come to the idea of the parametric effect.

Parametric amplification means production of inhomogeneities of the 'growing' mode which, in fact, means the creation of new perturbations (Lukash 1980). The idea is very simple. If only the 'growing' mode of perturbation were excited at the beginning, then it would be conserved until $\kappa \ll 1$ independently of the expansion rate (i.e. for any equation of state):

$$q_1 = C_1, \qquad \kappa \ll 1. \tag{5}$$

Therefore, only the presence of the 'decaying' mode at the initial stage can increase the 'growing' mode amplitude at a late stage. Indeed, let us suppose that the condition $a \sim \eta$ no longer holds for a short time: i.e. $a \sim \eta$ at $\eta < \eta_1$, $a = a(\eta)$ at $\eta_1 < \eta < \tilde{\eta}_1$, and again $a \sim \eta$ at $\eta > \tilde{\eta}_1$. Also let us assume the 'natural' random phase initial conditions ($\eta < \eta_1$):

$$C_1 \sim C_2 \ll \kappa \ll 1. \tag{6}$$

The point is that the thorough separation of perturbation on the modes (by the initial phase) is possible only at the first and the third stages, when $a \sim \eta$ (see equation 1). During the intermediate stage they mix and a part (one half roughly speaking) of the 'decaying' mode transforms into the 'growing' mode. Hence, at $\eta > \tilde{\eta}_1$ the mode amplitudes are

$$\tilde{C}_1 \approx C_1 + \frac{C_2}{2\kappa_1} \sim \frac{C_1}{\kappa_1} \gg C_1; \qquad \tilde{C}_2 \approx \tfrac{1}{2} C_2 \sim C_1, \tag{7}$$

i.e. we have very large amplification in the 'growing' mode and in the perturbation energy, while the total level of perturbation (the sum of the two modes) is practically unchanged.

The physical reason for this amplification is a non-stationary background. It reminds us of the creation of γ-quanta in the electromagnetic resonator; that is why we call this effect parametric (or non-adiabatic) creation. It is also very similar to the production of gravitons (see Grishchuk 1974) or a minimally coupled scalar field in the expanding universe, only in our case we deal with phonons, the sound quanta.

Resuming, we can say that the parametric effect provides exactly what we need: starting from the 'natural' initial conditions $C_1 \sim C_2 \ll \kappa_1$, we obtain the required ones $C_2 \ll C_1 < 1$. A desirable amplitude $C_1 \sim 10^{-4}$ depends on the seed initial perturbation amplitude and on the κ_1 parameter (the earlier η_1 the greater C_1).

We now discuss the realisation of this idea and the properties of the perturbations created. The scalar q satisfies the equation

$$q'' + \frac{2\alpha'}{\alpha} q + (\beta k)^2 q = 0, \tag{8}$$

where the prime denotes $\mathrm{d}/\mathrm{d}\eta$, $\alpha^2 = a^2(\epsilon + p)/(\mathrm{d}p/\mathrm{d}\ln\epsilon)$, $\beta^2 = \mathrm{d}p/\mathrm{d}\epsilon$ and $p = p(\epsilon)$. For $\kappa \ll 1$, the term $\sim k^2$ in equation (8) is negligible and the general solution is

$$q = C_1 + \overline{C}_2 \int \frac{\mathrm{d}\eta}{\alpha^2}, \tag{9}$$

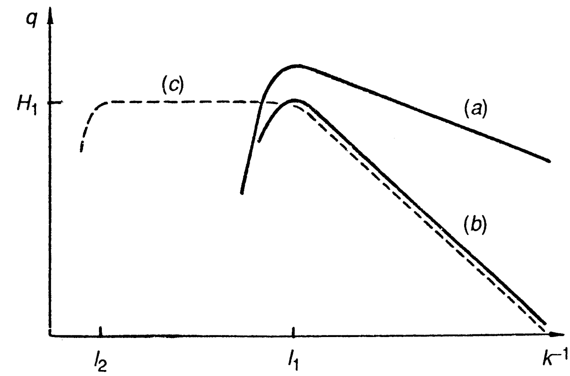

Fig. 1. Spectra of the primordial perturbations produced parametrically from the seed (*a*) thermal, or (*b*) quantum fluctuations of the gravitational potential ($H_1 \approx t_{Pl}/t_1$ and $l = a/a'$ where the scales $l_{1,2}$ coincide with the Hubble time at $\eta_{1,2}$ respectively) or (*c*) quantum fluctuations of a scalar (inflanton) field at the de Sitter intermediate stage ($H_1 = H_{dS}$).

where the constant $\overline{C}_2 = \overline{C}_2(k) \sim C_2$ in accordance with the initial conditions (1). The parametric amplification coefficient is then derived in a straight forward way (see Kompaneets *et al.* 1982). If the seed perturbations are gaussian (e.g. quantum vacuum point-zero fluctuations of gravitational potential), then the perturbations created will be gaussian as well with the spectra presented as in Fig. 1.

We will not go into further details (for a review see Lukash and Novikov 1985), but let us dwell upon two points. First, on the nature of the 'intermediate' state, modern cosmology proposes that this can be inflation. The cosmological model including the period of inflation, when the scale factor of the expansion is $a(t) \sim e^{Ht}$, at the beginning of the universe is currently the most popular. The inflationary scenario was born as a natural consequence of the application of grand unified gauge theories (GUTs) to the very early stages of the universe. On the other hand, some models of the inflationary universe have been proposed in which the cause of inflation is not GUT; for example, Starobinsky's theory (1980) of primordial inflation at the Planckian time or Linde's (1983) chaotic inflation scenario.

The inflationary universe scenario itself has solved many fundamental problems of cosmology and it is too fascinating to be abandoned. Thus, it seems natural to separate the inflation from the GUTs and from any specific scheme and analyse the different possibilities.

In concluding this Section, a few words about the place of the parametric effect among the different schemes of perturbation production. Practically, all creation effects have a parametric nature: all of them must produce the 'growing' mode on a scale much larger than the horizon scale. But various effects differ according to the different assumptions on the nature of the seed physical fields responsible for perturbations. Examples are quantum-gravitational one loop effects (Starobinsky 1980; Hartle and Horowitz 1981; Mukhanov and Chibisov

1981), gravitating scalar Higgs fields and vacuum phase transitions (Kirzhnits 1972; Guth 1981; Linde 1983; Bardeen *et al.* 1983), and quantum vacuum point-zero fluctuations of gravitational potential considered above, and so on.

Let us now turn to the further evolution of perturbations created near the singularity.

3. Further Evolution and Dark Matter

Work in recent years has showed that our universe is very likely to be multicomponent and to contain essentially weakly interacting particles. The physics of fundamental interactions points out that, alongside known particles, other collision-less particles (the hypothetical 'inos', monopoles, primordial black holes, etc.) should also be present in the cosmic medium. Observational cosmology also provides evidence for the non-baryonic nature of dark matter in the universe. Massive collision-less particles must exist in inflationary models of the universe. In such a situation the question arises of how the collision-less particles, both relativistic and nonrelativistic, affect the perturbation dynamics from the very early universe up to the nonlinear epoch of galaxy formation.

Further on, we present an investigation of this problem for the multicomponent universe. The analysis shows that the most important influence on the perturbation evolution comes from the so-called 'equality' epoch, where the densities of all relativistic and all nonrelativistic particles are equal ($Z \sim 10^4$ for the standard model). In this respect, the most interesting is a model with two components of weakly interacting particles—light and heavy—so that the relation between their particle numbers determines the 'equality' time (see Lukash 1987; Kahniashvili *et al.* 1987).

Below we assume that the present density of matter is provided by heavy relic (r) particles ($m \gg 10$ eV) and all the other collision-less components (ν) are light ($m < 10$ eV). In this case the light particles in the pre-recombination time are relativistic and the dependence of their distribution function on the momentum modulus is not important (only the local total density of ν-particles enters the equations in the early universe).

For the interacting particles of that time, the main contribution comes from photons (the γ component) which are easily described as an ideal fluid with the equation of state $p = \epsilon/3$. Thus, the model has only two free parameters, the ratios of the numbers of light and heavy collision-less particles to the total number of all relativistic particles including photons:

$$\nu = n_\nu/(n_\nu + n_\gamma), \qquad r = n_r/(n_\nu + n_\gamma). \tag{10}$$

For the standard hot dark matter (HDM) model $\nu = 0 \cdot 3$ and $r = 0 \cdot 16$; while for the standard cold dark matter (CDM) model $\nu = 0 \cdot 4$ and $r = 0$. The normalisation of the scale factor is chosen in such a way that at late stages, when the r-particles become nonrelativistic,

$$a = \eta(1 + \eta). \tag{11}$$

For $\eta \gg 1$, the r-component (dark matter) dominates the expansion and the large structure formations are ruled by the perturbations in this leading component. At this stage matter perturbations in other components no longer gravitationally influence the development of the r-component inhomogeneities. Thus, our model is good enough for finding perturbations in the cold dark matter since the changes in photons (recombination and post-recombination periods), which break the applicability of the y-component equation of state, happen at $\eta > 1$. For this reason, we also neglect here the contribution of baryons which slightly change the equation of state $p = \epsilon/3$ just before hydrogen recombination.

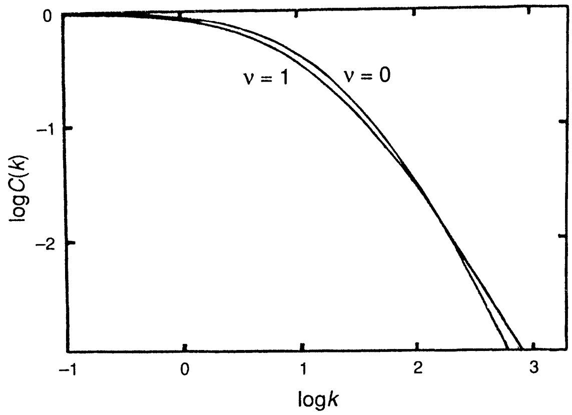

Fig. 2. Transfer functions of the adiabatic perturbations in the CDM model ($r = 0$) for $\nu = 0$ and $\nu = 1$.

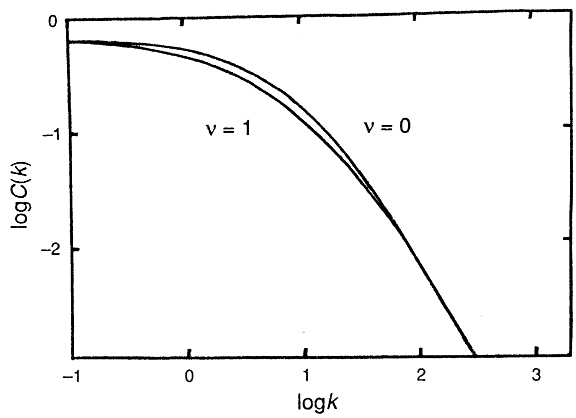

Fig. 3. Transfer functions of the isocurvature perturbations in the CDM model ($r = 0$) for $\nu = 0$ and $\nu = 1$.

4. Transfer Functions

The most convenient way to present the results of perturbation evolution is the transfer function $C(k)$. It is the ratio of the two spectra—the spectrum at a late stage just before the perturbations enter nonlinear evolution, and the primordial spectrum of perturbations created at the beginning of the expansion.

For the CDM model ($r = 0$), we show the transfer functions of the adiabatic (Fig. 2) and isocurvature (Fig. 3) perturbations for the limiting values of $\nu = 0$

and 1. All other curves for $0 < \nu < 1$ lie in between these two. Both demonstrate good scaling of the transfer functions: within an accuracy of ~10% they are independent of the ν-parameter. Such a weak dependence of $C(k)$ on ν is due to a successful choice of wavevector k normalisation (see equation 11). In the physical space all scales are strongly dependent on ν, but the shape of the spectrum of inhomogeneity [the $C(k)$ function] does not depend on ν.

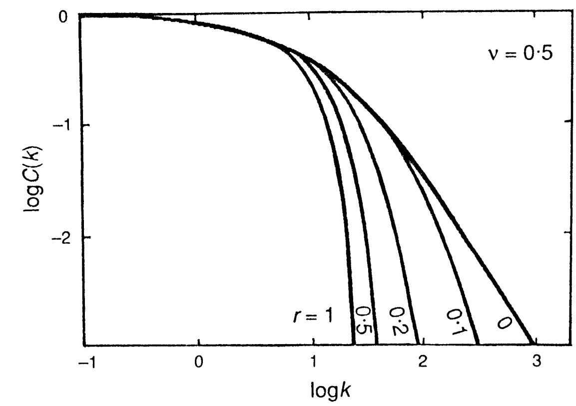

Fig. 4. Transfer functions of the adiabatic perturbations for different r values and $\nu = 0 \cdot 5$.

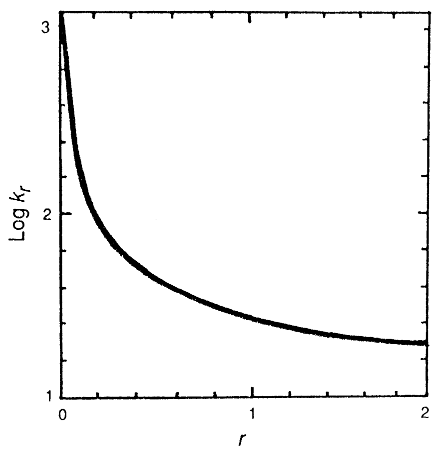

Fig. 5. Cutoff wavenumber k_r of the adiabatic transfer function $[C(k)_r = 10^{-3}]$ versus the r parameter.

Similar behaviour demonstrates a general model ($r \neq 0$). Here a new characteristic scale appears due to the free streaming of r-particles when they were relativistic: a cutoff where the transfer function sharply decays to smaller scales (see Fig. 4). The cutoff wavenumber k_r versus r is shown in Fig. 5.

The scaling effect takes place in the general case. For a fixed r, the transfer functions are ν invariant (within 10% independent of ν) in our special normalisation ($k = 1$ for the equality horizon).

5. Astrophysical Implications and Isocurvature Models

The essential independence of the transfer functions $C(k)$ on the ν-parameter (see figures) allows for several observational tests which can be easily done without a perturbation technique, but simply by using the background model equations. The first example is a characteristic cutoff scale in the perturbation spectrum which marks the change in the spectrum slope from the side of large scales in the CDM models. It coincides approximately with the transfer function shape change which takes place at $k_{eq} \sim 10$ with $C(k_{eq}) \sim 0 \cdot 5$ for adiabatic perturbations. In physical units this corresponds to

$$k = 10\lambda_{eq}/\lambda, \qquad \lambda_{eq} = 37h^{-2}(1-\nu)^{-\frac{1}{2}}\ \mathrm{Mpc}, \tag{12}$$

where λ is an up-to-date perturbation wavelength in Mpc. We see that the large structure scale λ_{eq} (about the supercluster or void scale) tends to infinity with $\nu \rightarrow 1$ (cf. the standard model $\lambda_{eq} \approx 45h^{-2}$ Mpc). In a similar way the spectrum cutoff scale and the amplitude values of the multipole $\Delta T/T(l)$ harmonics of the microwave background angular variations grow with increasing ν. Estimates for the ν-parameter upper limit without contradicting the relic data (Strukov *et al.* 1987) show that

$$\nu < 0 \cdot 8, \tag{13}$$

which is only twice as great as that in the standard model. Let us note that, contrary to other tests of relativistic particle species, our independent cosmological test evaluates the total number of all light ($m < 10$ eV) particles in the universe, including elusive particles (such as gravitons) whose discovery by other means now seems improbable.

Let us now turn to another interesting problem of the multicomponent universe: the possibility of a large perturbation in the non-dominating dark matter components. For the remainder of the discussion I follow an earlier paper (Lukash 1989).

Amongst different schemes for biased galaxy formation proposed recently there are a few that speculate on the idea of large baryonic perturbations: ranging from the hypothesis about the deficit of baryons in large cosmic voids embracing scales of 30–200 Mpc (cosmic hubbles) (Kofman *et al.* 1987), to the assumption than all the visible matter in the universe was born in a gigantic baryonic island extending to redshifts $Z \sim 4$ where baryons vanish beyond the cosmic islands (Kardashev *et al.* 1987; Dolgov *et al.* 1987). These and similar suggestions originate from a more general assertion that nondominating media do not manifest themselves dynamically now. It means that, even if they were highly perturbed now (but the dominating dark matter remains homogeneous), they would not contribute to the gravitational potential, i.e. to the 'growing' adiabatic mode of perturbations.

However, this idea misses one point: since these large perturbations were created very early (e.g. baryons were produced at GUT times $t \sim 10^{-35}$ s), they did influence the expansion dynamics during the equality epoch. As a result, they induced the 'growing' adiabatic mode with a high amplitude which, in turn, caused later large background $\Delta T/T$ fluctuations due to the Sachs–Wolfe effect and, as a consequence, contradictions with the observations.* Let us briefly outline the proof. We consider an early multicomponent universe consisting of the relativistic components (which include γ- and ν-particles—see Section 2) and of the nonrelativistic 'dust' r-particles. The latter include 'cold' dark matter and 'baryons' (b) whose portion is fixed by the parameter

$$\Omega_b = \Omega_b(\boldsymbol{x}) = \rho_b/\rho_r . \tag{14}$$

For simplicity, we neglect the baryonic pressure and therefore baryons and cold particles move together.† Thus, Ω_b is the integral of motion, i.e. it depends only on space coordinates in the synchronous comoving frame.

Now, we consider early evolution when the cosmological horizon was much less than a characteristic scale of variation of the Ω_b function:

$$t \ll l = |\,\Omega_b/\nabla\Omega_b\,| . \tag{15}$$

Let all the components be initially at rest ($t \to 0$). Then under the condition (15) the metric is locally isotropic,

$$ds^2 = dt^2 - (aR)^2(dx^2 + dy^2 + dz^2), \tag{16}$$

where $a = a(t, \boldsymbol{x})$, $R = \text{const.} \sim t_{eq}$,‡ and the only nontrivial equations are

$$\begin{aligned} \epsilon &= 3C_1^2/R^2a^4 , \qquad \rho_r = 12C_1\, C/R^2a^3 , \\ \overset{\bullet}{a}^2 &= C_1^2 + 4aC_1\, C, \quad (\bullet) = aR\,\partial/\partial t, \end{aligned} \tag{17}$$

where $C_{(1)} = C_{(1)}\,(\boldsymbol{x})$ are space functions fixed by initial conditions. The solution for the scale factor is

$$\begin{aligned} a &= C_1\, \tau(1 + C\tau), \\ t &= R\int a\, d\tau = \tfrac{1}{2}RC_1\, \tau^2(1 + \tfrac{2}{3}C\tau) . \end{aligned} \tag{18}$$

Let us now set the baryon excess to be extremely perturbed initially while the CDM and the dominating relativistic component are spatially homogeneous (for the physical mechanisms see Dolgov 1987). This means the following choice of $C_{(1)}$ functions (see equation 14):

$$C_1 = 1, \qquad C = 1/(1 - \Omega_b), \tag{19}$$

so that for the equality time

$$a_{eq} = 0\cdot 25(1 - \Omega_b), \qquad \tau_{eq} = 0\cdot 2(1 - \Omega_b) . \tag{20}$$

* In principle a special geometry (e.g. a high degree of spherical symmetry of the cosmic void or island or a certain position of the observer etc.) could conceal these temperature fluctuations from the observer, but here we consider an arbitrary density configuration.

† This simplification allows one to generalise the problem: instead of 'baryons' one may consider any heavy particles, e.g. the CDM particles which invert the problem. (Note that Ω takes any value from zero to unity.)

‡ The a function and R parameter are gauge invariant since the reference system (16) is fixed unambiguously by the condition $u_r^i = (1, 0, 0, 0)$.

For further estimates we substitute $\delta\Omega_b \approx \Omega_b \approx 0\cdot 1$ for the baryon perturbation on scales of $\sim l$.

At $\tau \ll \tau_{eq}$ when relativistic particles predominate, we have

$$a = (2t/R)^{\frac{1}{2}}, \qquad \epsilon = 3/4t^2,$$
$$\rho_r = \frac{6}{1-\Omega_b}(2R)^{-\frac{1}{2}}t^{-\frac{3}{2}} \ll \epsilon, \tag{21}$$

in accordance with the initial conditions. At $\tau \gg \tau_{eq}$ equations (17) and (18) yield

$$a = (1-\Omega_b)^{-\frac{1}{3}}(3t/R)^{\frac{2}{3}}, \qquad \rho_r = 4/3t^2,$$
$$\rho_b = 4\Omega_b/3t^2, \qquad \epsilon \ll \rho_r, \tag{22}$$

which, in fact, is a sum of two perturbation modes. The first line of (22) presents the first expansion term over the parameter $(t/l)^2 \ll 1$ of the 'growing' adiabatic mode [the quasi-isotropic solution of Lifshitz and Khalatnikov (1963)], while the second line describes the isocurvature perturbations in the baryons and in the non-dominating relativistic component.

Equations (21) and (22) display an important conclusion: gravitational field perturbations (scale factor perturbations here), vanished initially, arise in the nonrelativistic matter dominating era:

$$\delta a/a \approx \tfrac{1}{3}\delta\Omega_b \approx 0\cdot 03.$$

These large perturbations of the metric bring about cosmic microwave background fluctuations $\Delta T/T \sim 10^{-3}$–10^{-2} at angular scales $\sim 1°$–$90°$ which depend on the value of the variation scale of the baryon density perturbation l.

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