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Preprint No 65

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IN THE INFLATIONARY UNIVERSE

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Moscow - February 1985

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ABSTRACT

It is shown that not only adiabatic perturbations of density but also isothermal density perturbations are generated in the inflationary universe. The magnitude of the isothermal perturbations can be sufficiently large for the subsequent galaxy formation. The spectrum of isothermal perturbations in some models is different from the flat one.

1. The problem of generation of density perturbations which are necessary for galaxy formation is one of the most difficult cosmological problems. Fortunately, three years ago a simple and natural solution of this problem was suggested in the context of the inflationary universe scenario (for a detailed discussion of this scenario see ref. /1/). It was shown, that during inflation perturbations of density are generated with the spectrum $\frac{\delta \rho(\ell)}{\rho}$ which is practically independent of the length scale ℓ /2-7/ (the so-called "flat" spectrum, or Zeldovich spectrum /8/). In some theories the amplitude of density perturbations has the magnitude $\frac{\delta \rho}{\rho} \sim 10^{-4}$, which is just necessary for the galaxy formation /1/.

Unfortunately, the first success has led to some absolutization of the results obtained. It is commonly believed now, that only adiabatic perturbations with a sufficiently large magnitude can be generated during inflation. It is assumed that all other types of density perturbations (e.g. isothermal ones) always are negligibly small as compared with the adiabatic perturbations (see e.g. /9/). It is believed also that the spectrum of density perturbations in the inflationary universe is exactly scale independent or very slowly (logarithmically) grows with the growth of the wavelength ℓ . All these conclusions are actually valid for the simplest models considered in /2-7/. However, these conclusions together with the statement, that the energy density of the universe ρ is equal to the critical density ρ_c ($\Omega = \frac{\rho}{\rho_c} = 1$), may form something like the Procrustean bed for the theory of

formation of the large-scale structure of the universe, see e.g. /10/. Another problem is related to the experimental constraint $\frac{\Delta T}{T} \lesssim 5 \cdot 10^{-5}$ on fluctuations of temperature of the microwave background radiation. The theory of galaxy formation from the adiabatic perturbations with a flat spectrum predicts $\frac{\Delta T}{T} \sim \frac{1}{3} \frac{\delta \rho}{\rho} \sim 3 \cdot 10^{-5}$. To be more precise, this theory gives slightly different values of $\frac{\Delta T}{T}$ for different angles of observation θ (different wavelengths ℓ) and for different models of galaxy formation (neutrino-dominated universe, axion-dominated universe etc.) /11, 12/. At the present stage of the development of the theory, the results are also considerably author-dependent, but many authors think that the predictions for the quadrupole anisotropy of the background radiation (corresponding to the density perturbations of the wavelength ℓ of the same order as the size of the observable part of the universe, $\ell \sim 10^{28}$ cm) are now on the verge of contradiction with the experimental data (at least for the neutrino-dominated universe /11, 12/). To avoid this difficulty it would be better to have a spectrum of density perturbations which decreases with an increase of ℓ at 10^{22} cm $\lesssim \ell \lesssim 10^{28}$ cm.

Though many scientists do not consider these difficulties very seriously due to the large uncertainty both in experimental data and in theoretical estimates based on the theory of galaxy formation, it would be unwise to forget about these problems and not to study their possible solutions, since it may shed a new light on the problem of formation of the large-scale structure of the universe.

First of all one should try to understand whether the assumptions that $Q = 1$ and that the density perturbations are adiabatic ones with a flat spectrum are actually necessary.

The answer is that it is possible to obtain $Q \neq 1$ after inflation, but it is difficult to do it naturally, since it would require a fine-tuning of the degree of inflation /1/. Much easier is to change the shape of the spectrum of adiabatic perturbations by a proper choice of the effective potential $V(\varphi)$ of the inflaton (of the scalar field φ which governs inflation). However, it is not very easy to suggest a natural and realistic theory in which adiabatic perturbations would decrease at large ℓ . One more possibility is related to the isothermal perturbations. As it was shown in our recent paper /13/, such perturbations are also generated in the inflationary universe. At later stages of the universe evolution the isothermal perturbations may become dominating. Their magnitude can be sufficient for galaxy formation, whereas their spectrum in some models differs considerably from the flat one: Typically $\frac{\delta \rho(\ell)}{\rho}$ decreases at large ℓ /13/. A more detailed discussion of this question will be contained in the present publication.

2. First of all, let us remind, in a somewhat simplified way, the mechanism of generation of density perturbations in the inflationary universe /3-7/. Let us consider the universe with a scale factor $a(t) \sim e^{H(t-t_0)}$, containing a scalar field φ with a small effective mass, $m^2(\varphi) \equiv \frac{d^2 V}{d\varphi^2} \ll H^2$. In this case long-wave quantum fluctuations of the field φ are generated with the magnitude which at small

$t - t_0$ is given by /14, 15, 4/ ,

$$\langle (\delta\varphi)^2 \rangle = \frac{H^3}{4\pi^2} (t - t_0). \quad (1)$$

At large t $\langle (\delta\varphi)^2 \rangle$ asymptotically approaches /16/

$$\langle (\delta\varphi)^2 \rangle = \frac{3H^4}{8\pi^2 m^2}. \quad (2)$$

The contribution to the quantity $\langle (\delta\varphi)^2 \rangle$ from the fluctuations with momenta $\kappa_0 < \kappa < H$ is given by /1/

$$\langle (\delta\varphi)^2 \rangle_{\kappa_0} = \frac{H^2}{2\pi^2} \int_{\kappa_0}^H \left(\frac{\kappa}{H} \right)^{\frac{2m^2}{3H^2}} d\ln\kappa \quad (3)$$

This means that the magnitude of fluctuations of the field φ (at a logarithmic scale of κ) is given by

$$\delta\varphi(\kappa) \sim \frac{H}{\sqrt{2}\pi} \cdot \left(\frac{\kappa}{H} \right)^{\frac{m^2}{3H^2}}. \quad (4)$$

This leads to density perturbations of the field φ

$\delta\varphi(\kappa) \sim \frac{dV}{d\varphi} \delta\varphi(\kappa)$. For any $V(\varphi) \sim \varphi^n$ this yields

$$\frac{\delta\varphi(\kappa)}{\varphi} \sim \frac{\delta\varphi(\kappa)}{\varphi} \sim \frac{H}{\sqrt{2}\pi\varphi} \cdot \left(\frac{\kappa}{H} \right)^{\frac{m^2}{3H^2}}, \quad (5)$$

where φ is the average magnitude of the scalar field (see below). Note, that for $m^2 \ll H^2$ the spectrum $\frac{\delta\varphi(\kappa)}{\varphi}$ is almost exactly scale-independent /2-7/. For a more precise evaluation of the shape of this spectrum one should take into account the dynamical evolution of the field $\varphi(t)$ and the time-dependence of the factor $H(\varphi(t))$. For the adiabatic perturbations with $m \ll H$ it usually leads to a logarithmic dependence of $\frac{\delta\varphi(\kappa)}{\varphi}$ on κ /2-7/.

It will be important for us later that the main contribu-

tion to $\langle (\delta\varphi)^2 \rangle$ (2), (3) is given by fluctuations with exponentially large wavelengths,

$$\tilde{\ell} \sim \tilde{\kappa}^{-1} \sim H^{-1} \exp\left(\frac{3H^2}{m^2}\right). \quad (6)$$

This means that at scales smaller than $\tilde{\ell}$ the field φ changes very slowly. It looks as an almost homogeneous scalar field with an average amplitude

$$\tilde{\varphi} = \sqrt{\varphi_c^2 + \langle (\delta\varphi)^2 \rangle}. \quad (7)$$

Here φ_c is the classical part of the field φ , which becomes homogeneous after inflation /1/ and satisfies equation

$$\ddot{\varphi}_c + 3H\dot{\varphi}_c = -\frac{dV}{d\varphi_c} = -m^2\varphi_c. \quad (8)$$

Note, that at the scales $\ell \gg \tilde{\ell}$ (6) the magnitude of the field φ changes considerably and, for $\varphi_c = 0$, it may vary between $-O(\tilde{\varphi})$ and $+O(\tilde{\varphi})$.

3. A subsequent evolution of the density perturbations $\delta\rho_\varphi$ depends on the properties of the fields φ . In /3-7/ only one scalar field φ was considered, the inflaton field which governs inflation. In many cases it is sufficient for the study of density perturbations since the total energy density ρ_{tot} during inflation and soon after it is dominated by the inflaton energy ρ_φ , so that $\frac{\delta\rho_\varphi}{\rho_\varphi} \approx \frac{\delta\rho_{tot}}{\rho_{tot}}$. The corresponding density perturbations are adiabatic, and their amplitude is directly related to the amplitude of fluctuations of temperature of the ultrarelativistic gas created by the decay of the field φ : $\frac{\delta\rho_\varphi}{\rho_\varphi} \sim \frac{\Delta T}{T}$.

However, there exist many theories in which ρ_{tot} at the late stages of the universe evolution (when galaxies are formed) is dominated not by the inflatons and the products of their decay, but by particles (or classical fields) of a different type.

In unified theories of elementary particles there exist many scalar fields, some part of which (let us denote them by χ) extremely weakly interact with other particles (the so-called "hidden sector", see e.g. /17/). Such fields either are completely stable or they decay at some late stages of the universe evolution, and before their decay their energy density ρ_χ usually decreases much more slowly than the energy density of the hot ultrarelativistic gas $\rho_\varphi(T) \sim T^4$, which was formed due to the decay of inflatons φ in the process of reheating after inflation. Therefore the relative contribution of the fields χ to ρ_{tot} grows, and these particles under certain condition can give a dominating contribution to the present density of matter in the universe (hidden mass).

Let us consider as an example the field χ with the effective potential

$$V(\chi) = \frac{m^2}{2} \chi^2 + \frac{\lambda}{4} \chi^4, \quad (9)$$

and let us assume that the universe after inflation was reheated up to $T \gg \sqrt{m M_p}$ where $M_p \sim 10^{19}$ GeV is the Planck mass. The energy density of the hot dense gas created by the decaying inflaton field φ is given by

$$\rho_\varphi(T) = \frac{N \pi^2 T^4}{30} \approx \rho_{tot}. \quad (10)$$

where N is the effective number of particles with $m \ll T$; typically $N \sim 200$. The (effective) age of the hot universe is given by

$$t = \frac{1}{4\pi} \sqrt{\frac{45}{\pi N}} \frac{M_P}{T^2}. \quad (11)$$

Classical field χ satisfies equation analogous to eq. (8):

$$\ddot{\chi} + \frac{3}{2t} \dot{\chi} = - \frac{dV}{d\chi} = -m^2 \chi - \lambda \chi^3. \quad (12)$$

Due to the "friction" term $\frac{3}{2t} \dot{\chi}$, the field χ becomes convergently oscillating near $\chi=0$ only at $m^2 + 3\lambda \chi^2 \gtrsim t^{-2}$. In the case $m^2/\lambda \ll \chi_0^2 \ll M_P^2$, where χ_0 is an initial value of the field χ , from eqs. (10), (11) it follows that the oscillations start at

$$\frac{S_{tot}}{M_P^2} \gtrsim \lambda \chi_0^2, \quad (13)$$

i.e. at

$$S_\chi \sim S_{tot} \cdot \frac{\chi_0^2}{M_P^2}, \quad T^2 = T_0^2 \sim \sqrt{\frac{30\lambda}{N}} \frac{\chi M_P}{\pi}. \quad (13')$$

At $\chi \gtrsim \frac{m}{\sqrt{\lambda}}$ the amplitude of the oscillating field χ decreases as a^{-1} with the expansion of the universe, see e.g. /18/. Consequently, $S_\chi \sim V(\chi) \sim \lambda \chi^4 \sim a^{-4}$. When χ becomes smaller than $m/\sqrt{\lambda}$, the energy of the field χ becomes decreasing more slowly, $S_\chi \sim V(\chi) \sim n^2 \chi^2 \sim a^{-3}$ /18/. This regime starts at the temperature T_* , where

$$T_*^2 \sim T_0^2 \cdot \frac{\chi_*^2}{\chi_0^2} \sim \sqrt{\frac{30}{\lambda N}} \frac{M_P m^2}{\pi \chi_0}. \quad (14)$$

With a further decrease of temperature, consequently

$$\frac{\delta \chi}{\delta \varphi} \sim \frac{\chi_0^2}{M_p^2} \frac{T_*}{T} \sim \frac{m}{T} \left(\frac{\chi_0}{M_p} \right)^2 \cdot \left(\frac{30}{\lambda N \pi^2} \right)^{1/4}. \quad (15)$$

From eq. (15) it follows that for a sufficiently large initial value of the field $\chi = \chi_0$ the field χ at late stages of the universe evolution (small T) gives a leading contribution to δS_{tot} see e.g. /19-21/. In such a case fluctuations of the field χ , which were neglected in /3-7/, may give a leading contribution to $\frac{\delta S_{tot}}{S_{tot}}$ at the late stages of the universe evolution. Let us consider three different examples which will illustrate some features of this mechanism of generation of density perturbations.

4. Let us assume first that the field χ originally was equal to zero (e.g. due to high-temperature effects /22/). Inflation generates long-wave fluctuations with a typical amplitude, which, according to eqs. (2), (7), is given by

$$\tilde{\chi} = \sqrt{\langle (\delta \chi)^2 \rangle} = \sqrt{\frac{3}{2}} \frac{H^2}{2\pi m}. \quad (16)$$

This leads to the appearance of density perturbations

$$\frac{\delta \rho_\chi(k)}{\delta S_{tot}} \sim \frac{\delta \chi(k)}{\chi} \cdot \frac{\rho_\chi}{S_{tot}} \sim \frac{\rho_\chi}{S_{tot}} \cdot \frac{m}{\pi H} \cdot \left(\frac{k}{H} \right)^{\frac{m^2}{3H^2}}. \quad (17)$$

These density perturbations are generated even if there were no adiabatic perturbations $\delta \rho_\varphi$ and no perturbations of temperature $\delta T \sim T^{-3} \delta \rho_\varphi$. (In what follows we assume adiabatic perturbation to be small). Moreover, since the χ -particles are stable (or almost stable) and since originally $\rho_\chi \ll \rho_{tot}$, perturbations $\delta \rho_\chi$ do not induce any perturbations of temperature at the early stages of the

universe evolution. In this sense the perturbations δg_x (17) can be called isothermal. (Note, that these perturbations in general are not related to the isothermal perturbations of a relative baryon density n_B/n_Y).

As follows from eq. (15), the coefficient $\frac{\delta g_x}{\delta_{tot}} \sim \frac{\delta g_x}{\delta_x + \delta_y}$ initially is small, but it can asymptotically approach unity at late stages of the universe evolution. In such a case one can have $\frac{\delta g_x}{\delta_{tot}} \sim 10^{-4}$ for $\frac{m}{H} \sim 10^{-4}$, which is quite reasonable.

If the field χ decays early, so that the value δ_x/δ_{tot} always remains small (note that initially $\frac{\delta_x}{\delta_{tot}} \sim \frac{H^2}{M_P^2} \ll 1$), still one can obtain isothermal perturbations of necessary magnitude $\frac{\delta g_x}{\delta_{tot}} \sim 10^{-4}$ if $\delta_x/\delta_{tot} \sim 10^{-3}$ at present and if $m \sim cH$, where c is not much smaller than unity. This possibility seems to be rather natural and interesting. Indeed, there exist many theories in which scalar fields acquire an induced mass $m \sim cH$ during inflation. This occurs e.g. if the field χ interacts with the scalar curvature R as follows: $\Delta L \sim \frac{c}{12} \chi^2 R$. (Note, that even if such a term is absent from the classical Lagrangian, it may appear due to quantum corrections, with $c \sim e^2 \sim 10^{-1} / 23$). The same effect typically occurs also in the theories based on $N=1$ supergravity, where the fields χ in the hidden sector often acquire mass $\sim cH$ due to their induced interaction with the inflaton field φ /20, 21/. This implies that the spectrum (17) should not necessarily be very flat: In many cases mentioned above

$$\frac{\delta \rho_x(\kappa)}{\rho_{tot}} \sim \frac{\rho_x}{\rho_{tot}} \cdot \frac{c}{H} \left(\frac{\kappa}{H} \right)^{\frac{c^2}{3}}, \quad (18)$$

where the constant C^2 is model-dependent and may take, in particular, any value from $O(1)$ to $O(10^{-2})$. The theory of galaxy formation with such a spectrum of isothermal perturbations may differ considerably from the standard adiabatic theory (see e.g. /24- 26/). Note, that the amplitude of fluctuations $\frac{\delta \rho_x(\kappa)}{\rho_{tot}}$ decreases at large $\ell \sim \kappa^{-1}$, which is just desirable to overcome possible difficulties with the quadrupole anisotropy of the thermal background radiation.

It is worth noting also, that eq. (17) describes the average magnitude of fluctuations only. As we have mentioned already, the value of the field χ varies from $-O(\tilde{\chi})$ to $+O(\tilde{\chi})$ in different exponentially large domains of the universe /1/. Therefore the value of $\frac{\delta \rho_x}{\rho_x}$ in different domains of the universe may vary considerably, and it is quite possible that we live not in a typical domain with $\delta \rho_x / \rho_{tot}$ given by (17) but in some domain in which, $\frac{\delta \rho_x(\kappa)}{\rho_{tot}}$ is better fitted for creation of galaxies of our type.

5. In the previous example we assumed that the field χ initially was in the minimum of $V(\chi)$. However, as it is emphasized in /1, 27, 28/, this assumption is correct only if the field χ sufficiently strongly interacts with other fields (or if it has an induced mass $m_\chi \sim cH$, as discussed above) and if it is sufficiently small ab initio. In a general case one should study evolution of an initial distribution of the fields φ and χ according to the chaotic inflation scenario /27/. With this purpose let us

consider a theory of two noninteracting fields φ and χ with $V(\varphi) = \frac{m_\varphi^2}{2} \varphi^2 + \frac{\lambda_\varphi}{4} \varphi^4$, $V(\chi) = \frac{m_\chi^2}{2} \chi^2 + \frac{\lambda_\chi}{4} \chi^4$. Let us assume that $m_\chi^2 \gg m_\varphi^2$, $\lambda_\varphi \gg \lambda_\chi$, and that the field χ very weakly interacts with all other fields. As is shown in /27/, at $\varphi, \chi \gtrsim M_p$ these fields roll down to the minima of $V(\varphi)$, $V(\chi)$ very slowly due to the large friction terms of the type $3H\dot{\varphi}$ (8), and the universe with $\varphi, \chi \gtrsim M_p$ expands quasiexponentially. Typically, due to the greater curvature of the potential $V(\chi)$, the regime of slow rolling finishes first for the field χ . This occurs when the Hubble parameter $H(\varphi, \chi) = \sqrt{\frac{8\pi}{3M_p^2} (V(\varphi) + V(\chi))}$ becomes smaller than the effective mass $m(\chi) = \sqrt{\frac{d^2 V}{d\chi^2}}$, which at large χ is given by $\sqrt{\lambda_\chi} \chi$. Therefore usually the last stage of inflation is governed by the lightest of the fields φ and χ (in this respect see also /29/). Inflation finishes at $\varphi \sim \frac{M_p}{3}$, when $H \sim \sqrt{\lambda} M_p / 1, 27/$. In this case from eq. (2) we obtain the well-known result /1-7/

$$\frac{\delta \mathcal{S}_\varphi}{\mathcal{S}_\varphi} \sim \sqrt{\lambda_\varphi}. \quad (19)$$

Thus, the amplitude of adiabatic perturbations in this model is given by the square root of the smallest of the coupling constants λ_φ , λ_χ .

Due to the existence of the friction term $3H\dot{\chi}$, the field χ during inflation rolls down not to $\chi=0$, but to such a region, where

$$m^2(\chi) \sim \lambda \chi^2 \sim c^2 H^2(\varphi) = c^2 \cdot \frac{8\pi}{3M_p^2} V(\varphi), \quad (20)$$

where C is some constant, $C \lesssim 1$. Since at the last stages of inflation $\varphi \sim M_P$, eq. (20) together with (5) implies that simultaneously with the perturbations of $\mathcal{S}_\varphi$, perturbations of $\mathcal{S}_\chi$ are also generated,

$$\frac{\delta \mathcal{S}_\chi}{\mathcal{S}_\chi} \sim \sqrt{\lambda_\chi} . \quad (21)$$

Thus, the amplitude of isothermal perturbations in this model is given by the square root of the largest of the coupling constants $\lambda_\varphi, \lambda_\chi$.

It is interesting to note that in this example, just as in the example considered before, $m(\chi) \sim C H$. This leads again to the power-law spectrum $\delta \mathcal{S}_\chi(\varphi) \sim \sqrt{\lambda_\chi} \cdot \left(\frac{\kappa}{H}\right)^{\frac{1}{2}}$, which decreases at large distances $\ell \sim \kappa^{-1}$. This result is valid for a sufficiently slowly changing Hubble parameter $H(\varphi(t))$. For a rapidly decreasing $H(t)$ one should also take into account the decrease of the magnitude of fluctuations $\delta \varphi(\kappa) \sim H(t)$ with relatively large κ produced at the last stages of inflation. We wish to emphasize, that whereas the time dependence of $H(\varphi(t))$, the spectrum of $\delta \varphi(\kappa)$ and the spectrum of adiabatic perturbations $\delta \mathcal{S}_\varphi(\kappa)$ are determined by $V(\varphi)$ only, the spectrum of $\delta \chi(\kappa)$ at a given $H(\varphi)$ is determined by $V(\chi)$. This makes it rather easy to alter the shape of the spectrum of isothermal perturbations by a proper choice of $V(\varphi)$ and $V(\chi)$.

6. The last example to be considered in this paper is related to the axion-type models, in which the effective potential is periodic in χ , e.g. $V(\chi) = \frac{1}{2} V(\pi \chi_0) (1 - \cos \frac{\chi}{\chi_0})$. Isothermal perturbations in such models also can be large.

A rather interesting effect which occurs in such theories is a specific dumping of the large-scale perturbations of $\frac{\delta \rho}{\rho}$ due to the periodicity of $V(x)/13/$.

The curvature of $V(x)$ in the models of this type typically is extremely small, m_x is many orders smaller than H during inflation. Therefore fluctuations of the field x in this model are given by (3), (4) with

$$\delta x(k) \sim \frac{H}{\sqrt{2} \pi} \quad (22)$$

However, eq. (5) in this model holds only for density perturbations with sufficiently large k , with the total dispersion smaller than πx_0 . The reason is that if the total dispersion of fluctuations with $k_0 < k < H$ (3)

$$\tilde{x}(k_0) = \sqrt{\langle (\delta x)^2 \rangle_{k_0}} = \frac{H}{\sqrt{2} \pi} \ln^{\frac{1}{2}} \frac{H}{k_0} \quad (23)$$

becomes greater than πx_0 , then the field x becomes uniformly distributed at the circle $-\pi x_0 < x < \pi x_0$. This means that inside the domain of the universe of a size $\ell \gtrsim k_0^{-1}$ all values of the field x in the interval $-\pi x_0 < x < \pi x_0$ appear with the same probability, and the average energy density is given by $\frac{1}{2} V(\pi x_0)$. In this case any long-wave fluctuation $\delta x(k)$ with $k \ll k_0$, which is equivalent to a rotation of the distribution of the field x by an angle $\delta x(k)$, does not change the average distribution of $V(x)$ inside the domains of a size $\ell \sim k^{-1}$. Thus the perturbations of the field x with $k \ll k_0$ do not produce any density perturbations $\delta \rho_x(k)$.

The value of κ_0^{-1} will be of the same order as the size of the observable part of the universe $\ell \sim 10^{28}$ cm if $\kappa_0^{-1} \sim H^{-1} \exp(70)$. This yields $\chi_0 \sim 0.7 H$. For $\chi_0 \gg 0.7 H$ the spectrum of $\delta \rho_x(\kappa)$ will be flat at all scales of interest.

Now let us estimate the magnitude of isothermal density perturbations $\delta \rho_x / \rho_{tot}$. Let us assume first that $\chi_0 \sim 10^{12}$ GeV. In this case $\rho_x \approx \rho_{tot}$ at the late stages of the universe evolution /19/, and

$$\frac{\delta \rho_x}{\rho_{tot}} \sim \frac{\delta \rho_x}{\rho_x} \sim \frac{\delta \chi}{\chi} \sim \frac{H}{\chi}. \quad (24)$$

To get $\delta \rho_x / \rho_{tot} \sim 10^{-4}$ one should have $H \sim 10^8$ GeV at the last stages of inflation, and $\chi_0 \sim 10^{12}$ GeV. In this case $\chi_0 \gg H$ and the spectrum is almost exactly flat. If, however, $\chi_0 \ll 10^{12}$ GeV (say, $\chi_0 \sim 10^{10}$ GeV), then $\rho_x / \rho_{tot} \ll 1$, and one can get $\delta \rho_x / \rho_{tot} \sim 10^{-4}$ at $H \sim \chi_0$. In such a case the effect of damping of $\delta \rho_x(\kappa)$ at large $\ell = \kappa^{-1}$ may become important.

(Note, that from our results it follows that the "hidden particles" χ can play an important role in the process of galaxy formation even if their density ρ_χ is much smaller than the total energy density ρ_{tot} , since even in this case perturbations of ρ_χ may be very large, $\delta\rho_\chi \gtrsim \delta\rho_\varphi$).

The last comment we would like to make about the axions in this paper is related to the usual constraint $\chi_0 \lesssim 10^{12} \text{ GeV} / 19$. In the inflationary cosmology this constraint is not unavoidable (for a discussion of this question see also [30]). Indeed, as we have already discussed, after inflation the universe becomes divided into many exponentially large mini-universes with different values of the field χ inside each of them. Whereas a typical value of the field χ is given by χ_0 , there may exist many exponentially large domains with $|\chi| \ll \chi_0$, and in one of these mini-universes we may live now. (Note, that the

optimal conditions for the existence of life of our type may well exist at χ much different from χ_0 , since the initial value of χ determines the amplitude of density perturbations (24), the relative baryon abundance at late stages of the universe evolution etc.).

In this paper we considered three different models in which isothermal density perturbations can be generated. Many other models of this type can be suggested. For example, in /31/ it was pointed out that the classical axion field changes the degree of GP-violation and therefore it can change the effectiveness of the process of baryosynthesis. This could lead to inhomogeneities of the baryon density in the universe. However, since after inflation all classical fields become almost exactly homogeneous, it was unclear how the abovementioned mechanism could work. Now it is clear that the inhomogeneities of the field χ can be created during inflation by the mechanism discussed above. We expect, however, that due to the small efficiency of the standard mechanism of baryosynthesis, this mechanism of generation of isothermal perturbations will be subdominant as compared with the mechanism considered in the present paper.

A similar mechanism can work rather efficiently in supersymmetric $SU(5)$ model, in which the baryon asymmetry of the universe may arise due to the generation of long-wave fluctuations of classical squark-slepton fields during inflation /32, 33/. Isentropic component of density perturbations of a slightly different type arises after decay of χ -

particles at later stages of the universe evolution, since this process leads to a relative increase of n_B/n_Y in those parts of the universe in which the energy density of the field χ was small.

The main conclusion which follows from this discussion is that the adiabatic perturbations with a flat spectrum should not be considered as the only perturbations which appear after inflation and which can be used in the "cosmological engineering". The present investigation shows that in many different theories of elementary particles the isothermal perturbations of density with a sufficiently large magnitude and with a spectrum different from the flat one can be generated /13/. This makes possible a more flexible approach to the theory of formation of the large-scale structure of the universe and to the investigation of the anisotropy of the thermal background radiation.

I am grateful to G.V.Chibisov, A.G.Doroshkevich, M.Yu. Khlopov, L.A.Kofman, V.P.Mukhanov, I.D.Novikov and A.A.Starebinsky for many useful and stimulating discussions.

REFERENCES

1. A.D.Linde, Rep. Prog. Phys. 47 (1984) 925.
2. V.F.Mukhanov and G.V.Chibisov, JETP Lett. 33 (1981) 532.
3. S.Hawking, Phys. Lett. 115B (1982) 295.
4. A.A.Starobinsky, Phys. Lett. 117B (1982) 175.
5. A.H.Guth and S.-Y. Pi, Phys. Rev. Lett. 49 (1982) 1110.
6. J.Bardeen, P.Steinhardt and M.Turner, Phys. Rev. D27 (1983) 679.
7. R.Brandenberger, Rev. Mod. Phys., to be published.
8. Ya.B.Zeldovich, Mon. Not. R. Astron. Soc. 160 (1972) 1 p.
9. M.Axenides, R.Brandenberger and M.Turner, Phys. Lett. 126B (1983) 178.
10. M.Davis, G.Efstathiou, C.Frenk and S.M.White, Preprint NSF-ITP/84-129 (1984).
11. A.A.Starobinsky, Pisma Astron. Zh. 9 (1983) 579.
12. V.N.Lukash, H.D.Naselsky and I.D.Novikov, in Quantum Gravity (Proceedings of the 4-th seminar on quantum gravity, Moscow 1984) to be published by World Scientific.
13. A.D.Linde, Pisma ZhETF 40 (1984) 496.
14. A.Vilenkin and L.Ford, Phys. Rev. D26 (1982) 1231;
15. A.D.Linde, Phys. Lett. 116B (1982) 335;
16. T.Bunch and P.Davies, Proc. Roy. Soc. 360A (1978) 117.
17. H.P.Willes, Phys. Rep., to be published.
18. M.Turner, Phys. Rev. D28 (1983) 1243.
19. J.Freskill, M.B.Wise and P.Wilczek, Phys. Lett. 120B (1983) 127;
20. L.F.Abbott and P.Sikivie, Phys.Lett. 120 B (1983) 133;
21. M.Dine and W.Fischler, Phys. Lett. 120B (1983) 137.

20. G.D.Coughlan, W.Fischler, E.W.Kolb, S.Raby and G.G.Ross,
Phys. Lett. 12 131B (1983) 59.
21. A.S.Goncharov, A.D.Linde and M.I.Vysotsky, Phys. Lett.
147B (1984) 279.
22. D.A.Kirzhnits, JETP Lett. 15 (1972) 529;
A.D.Linde, Rep. Prog. Phys. 42 (1979) 389.
23. A.D.Linde, Phys. Lett. 114B (1982) 431.
24. P.J.E.Peebles, The large scale structure of the universe
(Princeton Univ. Press: Princeton 1980).
25. S.M.White and M.J.Rees, Mon. Not. R. Astr. Soc. 183
(1978) 341.
26. G.J.Hogan, Astroph. J. Lett. 284 (1984) L1.
27. A.D.Linde, Phys. Lett. 129 B (1983) 177.
28. A.D.Linde, Phys. Lett. 132B (1983) 317.
29. L.A.Kofman, A.D.Linde and A.A.Starobinsky, Prepr.
Inst. of Astrophys. Est. SSR (1985).
30. S.-Y.Pi, Phys. Rev. Lett. 52 (1984) 1725.
31. M.Yoshimura, Phys. Rev. Lett. 51 (1983) 439.
32. I.Affleck and M.Dine, Princeton Univ. preprint (1984).
33. A.D.Linde, in "Quantum Gravity" (Proceedings of the
4-th seminar on quantum gravity, Moscow 1984), to be
published by World Scientific;
A.D.Linde, The new mechanism of baryogenesis and
inflationary universe, Lebedev Phys. Inst. preprint (1985).
34. A.Vilenkin and A.E.Everett, Phys. Rev. Lett. 48 (1982)
1867.

Т - 06423

Подписано в печать 8 февраля 1985 года

Заказ № 122. Тираж 100 экз. п/л 1,0

Отпечатано в Отделе научно-технической
информации ФИАН СССР

Москва, В-312, Ленинский проспект, 53