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Antiproton Interactions with Light Elements as a Test of GUT Cosmology.

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1. – Introduction.

One of the most intriguing problems of modern physics is the problem of the unification of strong, electromagnetic and weak interactions in the framework of a unified theory. A number of models of this theory has been developed. The theory of the grand unification provides a description of the interaction between particles at superhigh energies, and the most nontrivial predictions of the existing GUT models concern energies of the order of $\sim (10^{14} \div 10^{15})$ GeV at which the unification of strong, electromagnetic and weak interactions is hoped to take place. A direct experimental check of GUT predictions in this energy range is obviously possible neither in any planned accelerators nor in cosmic-ray experiments. This is why the very few GUT predictions (the proton instability, the n-$\bar{\text{n}}$ oscillation, the existence of the neutrino mass, etc.) at energies achievable at present give rise to such a great interest of the physical community.

On the other hand, at the very early stages of the cosmological expansion of the Universe the GUT energies that were inevitably realized must have left some footprints on the whole successive evolution of the Universe. The explanation of the observed baryon asymmetry within the framework of GUT's scenarios of the evolution of the very early Universe is a well-known example of the influence of the GUT's effects on the fate of the Universe. There are some other possibilities for checking the predictions of GUT models based on their astrophysical implications. In this review we intend to discuss the possibility of the verification of GUT models based on the fact that some GUT models predict the existence of sources of antimatter in the Universe. Thus the rather old problem concerning the existence of antimatter in the Universe is revived in a nontrivial way on the basis of GUT cosmology. Almost all earlier approaches to this problem were based on baryon-symmetric cosmology. In our review the possibility of the existence of antimatter in the Universe provided by some GUT models is considered mainly within the frame of baryon-asymmetric models.

One of the most important and *testable* implications of the presence of antimatter in the Universe is its influence on the abundance of the light elements. The degree of such an influence depends both on the amount of antimatter and on the dynamics of antiproton-nucleus interactions. Unfortunately, the antiproton interaction with nuclei is very poorly studied. So, one of the aims of this review is to stress the importance of the study of $\bar{p}A$ interactions. Since ^4He is the most abundant (after hydrogen) element in the Universe, the investigation of the $\bar{p}\,^4He$ interaction is of special interest. An important experiment of this kind already planned at LEAR which may provide useful information for astrophysics is the experiment PS-179 where an helium streamer chamber is used (Dubna-Frascati-Padova-Pavia-Torino Collaboration). It will be shown in this review that the results of these experiments may provide valuable information for putting restrictions on the parameters of some GUT models.

The review is organized as follows. We shall start with a brief description of the modern views on the evolution of the Universe (sect. **2**). After that we discuss the problem of the existence of antimatter in the Universe and its possible astrophysical observational effects (sect. **3**). We will show that unique information on the possible presence of antimatter at the stage of radiation dominance in the Universe ($t < 10^{13}$ s) may be obtained from the study of $\bar{p}\,^4He$ annihilation (sect. **4**). The difficulties of early attempts to include antimatter in cosmological considerations will be described in sect. **5**. The sources of antimatter, predicted by some GUTs, are discussed in sect. **6-8**. We consider the physics of primordial black holes (their formation and evaporation), the formation and annihilation of antimatter domains, the relationship between these sources of antimatter and the phase transitions in the early Universe as well as heavy metastable particles predicted by GUTs. Some aspects of the theory of galaxy formation are in close touch with these problems. In sect. **9** the observed abundances of light elements and their possible changes due to annihilation are discussed. In conclusion (sect. **10**) the necessary measurements and the restrictions on the parameters of GUT models which may be obtained from these measurements are considered.

The review cannot at any rate be treated as a complete or detailed description of the big-bang theory or grand unified models. There is no need of this in view of existing books and reviews [1-10]. Only those aspects of these theories which are in touch with the aims declared are discussed. However, to show the experimentalists in nuclear physics, to whom this review is addressed, the ropes of the modern cosmology, we include in this review some well-known basic points of the big-bang theory.

2. – The big-bang Universe.

2.1. *The main parameters of big-bang cosmology.* – The modern theory of the expanding Universe is based on the assumptions of its homogeneity and isotropy on large scales.

The expansion of the Universe proceeds in accordance with the Hubble law

$$V = Hr\,, \tag{2.1}$$

where r is the distance between the objects, V is the speed of expansion and H is Hubble's constant. The values of the modern magnitude of Hubble's constant, obtained in astronomical observations, are in the interval $(50 \div 100)$ km/s$\cdot$Mpc [11]. The magnitude of Hubble's constant determines the magnitude of the so-called « critical density »

$$\varrho_c = \frac{3H^2}{8\pi G}\,, \tag{2.2}$$

where G is the gravitational constant.

For $H = 50$ km/s$\cdot$Mpc the value of ϱ_c is

$$\varrho_c = 5\cdot 10^{-30}\ \text{g/cm}^3\,.$$

If the cosmological density ϱ (the density averaged over space regions greater than 100 Mpc) exceeds the critical one, the world is closed and the observed expansion will be inevitably succeeded by compression. The dimensionless quantity $\Omega = \varrho/\varrho_c$ is usually introduced, so that $\Omega > 1$ corresponds to the closed world. Values $\Omega < 1$ correspond to the open world; in this case the presently observed expansion will never stop.

The observational data on the value of ϱ (and, consequently, Ω) are ambiguous (see [12, 13]). Estimates of the mean density of matter in galaxies and clusters of galaxies based on their luminosity give the value $\varrho \sim 1.5 \cdot 10^{-31}$ g/cm³. On the other hand, the estimates of the masses of clusters of galaxies, from the velocity distribution of galaxies in the cluster and with the use of the theorem of the virial, turn out to be an order of magnitude higher than the estimates of the mass of the luminous matter, *i.e.* the so-called « hidden mass » paradox takes place—the largest part of matter must be hidden in the form of nonluminous objects. We shall return to some aspects of this paradox in sect. **9**.

The most important proof of the big-bang theory was the discovery of the isotropic microwave radio background—the relic radiation. Its observed flux F_ν has a thermal spectrum

$$F_\nu = \frac{2h\nu^3}{c^2} \frac{1}{\exp[h\nu/kT] - 1}, \tag{2.3}$$

where T is the temperature of the background radiation and ν is its frequency.

Calculations of the temperature of the background radiation based on the flux observed in different directions from the sky give for fixed ν the same value for the temperature T with an accuracy of less than 10^{-4} [10]. But for different ν, the measured value of T is much worse and lies within the interval $(2.65 \div 2.9)$ K. These uncertainties in T either may be experimental or they may reflect a real deviation of the background radiation from the thermal distribution. All our subsequent considerations will be based on the assumption of the thermal character of the background radiation. We will discuss, however, some possible physical processes leading to a distortion of the background spectrum.

The presence of a thermal electromagnetic background with the quoted temperature corresponds to the following mean number of relic photons per unit volume:

$$n_\gamma = 20T^3 = (350 \div 500)\ \text{cm}^{-3}. \tag{2.4}$$

The mean number of baryons per unit volume is

$$n_{\text{B}} \sim \varrho_{\text{B}}/m_{\text{p}} \sim (10^{-6} \div 10^{-7})\ \text{cm}^{-3}, \tag{2.5}$$

where m_{p} is the mass of the proton.

Therefore, the modern Universe is characterized by a very small baryon-to-photon ratio:

$$r_{\text{B}} = \frac{n_{\text{B}}}{n_\gamma} = 10^{-8} \div 10^{-10}. \tag{2.6}$$

The given interval of values of r_B accounts for the ambiguities in the magnitudes of the baryon density n_B and of the relic-radiation temperature T.

The thermal character of the background radiation is not changed by the cosmological expansion. Going back to the past, we obtain much higher temperatures. Thus extrapolation to the past leads to the picture of matter being in the state of hot plasma in equilibrium with radiation. The principal features of such a system are rather simple being governed by the well-known laws of thermodynamics. That is the bedrock of big-bang cosmology.

2·2. *The main stages of the evolution of the big-bang Universe.* – A brief sketch of the big-bang scenario will be given here.

Everything started from the « moment zero », when the expansion of the Universe began. Some ideas [14-17] on the physical nature of the beginning have appeared recently. But a quantitative theory of the beginning has not been constructed at present.

The time scale of the big-bang is depicted in fig. 2.1. The first character-

The beginning					VERY EARLY UNIVERSE
T_{Pl}	10^{19} GeV	t_{Pl}	10^{-43} s		
T_{GUT}	10^{15} GeV	t_{GUT}	10^{-35} s	GUT phase transition baryosynthesis: γ, ν, $\bar{\nu}$, W, Z, L, $\bar{\mathrm{L}}$, quarks, antiquarks, gluons	
T_{WS}	300 GeV	t_{WS}	10^{-11} s	WS phase transition	
T_{QCD}	300 MeV	t_{QCD}	10^{-5} s	QCD confinement $\mathcal{N}\bar{\mathcal{N}}$ annihilation: γ, ν, $\bar{\nu}$, e^+, e^-, π^0, $\mathcal{N}$, $\bar{\mathcal{N}}$	
		t	10^{-3} s	ν, $\bar{\nu}$, γ, e^+, e^-, n, p	
		t	1 s	n/p freezing	RD STAGE
		t	10^2 s	e^+e^- annihilation nucleosynthesis: γ, ν, $\bar{\nu}$, p, ^{3}He, d, ^{4}He, e^-	
		t	10^3 s		
		t_{RD}			MATTER DOMINANCE
		t_{rec}	10^{13} s	recombination, the growth of inhomogenieties: γ, ν, $\bar{\nu}$, neutral atoms	
		t	$(10^{16} \div 10^{17})$ s	Galaxy formation	
		now			

Fig. 2.1.

istic time is the Planck time

$$t_{\mathrm{Pl}} = \sqrt{\frac{G\hbar}{c^5}} = 10^{-43}\ \mathrm{s}\,, \tag{2.7}$$

where G is the gravitational constant.

In this time the size of the cosmological horizon, *i.e.* the distance travelled by a light signal during the given cosmological time, is equal to the Planck length

$$l_{\mathrm{Pl}} = \sqrt{\frac{G\hbar}{c^3}} = 3\cdot 10^{-33}\ \mathrm{cm}\,. \tag{2.8}$$

The values of l_{Pl} and t_{Pl} determine the space and time scales on which quantization of space-time is inevitable. Since quantum gravity theory has not been developed up to now, l_{Pl} and t_{Pl} are the boundaries of our present knowledge, *i.e.* we may evolve the theory of phenomena only for $t > t_{\mathrm{Pl}}$ and $l > l_{\mathrm{Pl}}$.

At the moment of $t \sim t_{\mathrm{Pl}}$ the temperature of the Universe was as high as

$$T \sim T_{\mathrm{Pl}} \sim m_{\mathrm{Pl}} c^2 \sim \sqrt{\frac{\hbar c^5}{G}} = 10^{19}\ \mathrm{GeV}\,.$$

Here $m_{\mathrm{Pl}} = 10^{-5}$ g is the Planck mass.

The period $t_{\mathrm{Pl}} \leqslant t \leqslant 1$ s refers to the very early Universe. Within the framework of GUTs it is possible, in principle, to answer the question on the physical conditions during this period. These conditions cannot be established unambiguously, since the parameters of GUTs as well as the correct GUT model itself have not yet been established. However, the standard prejudice of the not so far past of treating this period of expansion as a white spot of our knowledge is removed. The history of the first second may be pictured, and for a definite choice of GUT and its parameters the picture is definite. Such a close relationship between particle physics and the cosmology of the very early Universe provides a decisive check of some predictions of grand unified models by their cosmological consequences.

Existing or achievable in the near future experimental data from accelerators support the theoretical description of physical processes at temperatures below ~ 100 GeV. Cosmic-ray data, however ambiguous, may give some information on the physics below $T \sim 10^6$ GeV. Physical processes at the higher temperatures $T \geqslant 10^6$ GeV may be discussed only within the framework of GUTs.

The general tendency of the evolution of the very early Universe is given by the law of expansion of the Friedman Universe [1-5]

$$\varrho = \frac{3}{32\pi G t^2} = \frac{4.5\cdot 10^5}{t^2}\ \mathrm{g/cm^3} = \frac{3}{32\pi}\,\frac{m_{\mathrm{Pl}}^2}{t^2} \quad (\text{in units } \hbar = c = 1)\,, \tag{2.9}$$

where t is the cosmological time. At very high temperatures all the known particles (and the particles predicted by GUTs) were relativistic. They were in equilibrium with radiation, and the thermodynamics of a relativistic gas gives for the energy density ε the following relationship with the temperature T:

$$\varepsilon = \varrho c^2 = \varkappa \sigma T^4 , \tag{2.10}$$

where

$$\sigma = \frac{\pi^2 k^4}{15\hbar^3 c^3} = 7.57 \cdot 10^{-15}\ \mathrm{erg/cm^3 \cdot degrees^4} ,$$

$k = 1.38 \cdot 10^{-16}$ erg/degrees is the Boltzmann constant, and $\varkappa$ is the number of species of relativistic particles (taking into account their statistical weight). One obtains from (2.9) and (2.10) that

$$(kT)^4 = \frac{1}{\varkappa}\,\frac{45}{32\pi^3}\,\frac{\hbar^3 c^5}{Gt^2} \tag{2.11}$$

and

$$T = \frac{1.3\ \mathrm{MeV}}{(t/1\mathrm{s})^{\frac12}\varkappa^{\frac14}} = \frac{1.5 \cdot 10^{10}\ \mathrm{K}}{(t/1\mathrm{s})^{\frac12}\varkappa^{\frac14}} = \left(\frac{45}{32\pi^3}\right)^{\frac14} \frac{1}{\varkappa^{\frac14}} \sqrt{\frac{m_{\mathrm{Pl}}}{t}} \tag{2.12}$$

(in units $\hbar = c = k = 1$)

or

$$t = \frac{1.7\ \mathrm{s}}{(T/1\ \mathrm{MeV})^2 \varkappa^{\frac12}} = \frac{2.25 \cdot 10^{20}\ \mathrm{s}}{(T/1\mathrm{K})^2 \varkappa^{\frac12}} = \left(\frac{45}{32\pi^3}\right)^{\frac12} \frac{1}{\varkappa^{\frac12}}\,\frac{m_{\mathrm{Pl}}}{T^2} \tag{2.13}$$

(in units $\hbar = c = k = 1$) .

For the density of relativistic particles n_{r} (of all the kinds), we have (taking $3kT$ as the mean energy of the particles)

$$n_{\mathrm{r}} = \frac{\varepsilon}{3kT} = t^{-\frac32}\varkappa^{\frac14} \cdot 0.01 \left(\frac{c}{\hbar G}\right)^{\frac34} = \frac{\varkappa^{\frac14} \cdot 5 \cdot 10^{31}\ \mathrm{cm^{-3}}}{(t/1\mathrm{s})^{\frac32}} . \tag{2.14}$$

In any system, thermodynamical equilibrium is maintained, if the rate of the processes setting up the equilibrium is greater than the rate of change of the parameters of the system (its density, temperature, etc.). In the expanding Universe the latter practically coincides with the rate of expansion. When the time scale of the process exceeds the cosmological time scale (*i.e.* the time from the beginning of expansion), the equilibrium is broken. For particles with concentration n and relative velocities v the process with cross-section σ and rate σv has time scale τ

$$\tau = (n\sigma v)^{-1} , \tag{2.15}$$

so equilibrium with respect to this process is mantained in the expanding Universe at the cosmological time t, if the condition

$$\tau < t \tag{2.16}$$

is fulfilled. If the equilibrium condition (2.16) is broken for some particles, the phenomenon of « freezing-out » or « decoupling » takes place. The concentration of particles is no longer equal to the equilibrium one, their relative concentration with respect to the other particles is « frozen out », remaining unchanged at all the successive stages of the expansion.

The preceding discussion of the very early Universe was based on the simple consideration of the law of cosmological expansion (2.9). Within the framework of GUTs some specific features arise.

According to GUT, at very high temperatures, after the Planck time t_{Pl}, all the particles were massless. They acquired their mass as a result of transitions to the phase of spontaneously broken gauge symmetry, which must have occurred in the Universe when the temperature dropped below certain critical values. Even in the simplest version of GUT, based on the SU_5 gauge symmetry, two such transitions are predicted (see fig. 2.1) [7, 18].

1) At $T_{\text{GUT}} \sim 10^{15}$ GeV, *i.e.* at $t_{\text{GUT}} \sim 10^{-35}$ s, a phase transition from the SU_5 symmetric phase to the $SU_3 \times SU_2 \times U_1$ phase takes place. After this transition a difference between the strong and unified electroweak interactions arises. Before this phase transition transformation of all species of particles into any other species was possible. After the phase transition gauge bosons mediating baryon-number-nonconserving transitions turn out to be very massive $M_{\text{X}} \sim 10^{14}$ GeV and the respective processes soon become highly suppressed and go out of equilibrium. If there has been any excess of baryon number before t_{GUT}, it would have been removed due to baryon-number-nonconserving processes. But CP violation in these processes after t_{GUT}, when the equilibrium condition (2.16) is broken for these processes, induces baryon excess generation, and this baryon excess survives due to the high suppression of baryon number nonconservation at the later stages of the expansion. We discuss in some detail these processes in sect. **7**.

2) At $T_{\text{WS}} \sim 300$ GeV, *i.e.* at $t_{\text{WS}} \sim 10^{-11}$ s, a phase transition from the $SU_2 \times U_1$ symmetrical phase to the $U_{1,\text{e.m.}}$ phase of short-range (W and Z bosons acquire mass) weak interactions takes place. The masses of all the known particles are believed to be generated as a result of this phase transition.

The period between t_{GUT} and t_{WS} may show a rather complicated behaviour, determined by the parameters of GUT models. Other phase transitions, heavy-metastable-particle dominance stages, anomalous vacuum dominance stages, etc., may arise in this period in some versions of GUTs. We shall discuss these possibilities as well as their possible check in sect. **6-8**.

In the period after t_{WS} it may be expected that all the known types of quarks and antiquarks (as well as gluons) and leptons were in equilibrium. As the temperature dropped below the mass of heavy species, these particles and the corresponding antiparticles annihilated (or decayed), so that the corresponding contributions to the energy density were distributed between the lighter species.

The next characteristic moment (see fig. 2.1) is $t_{QCD} \sim 10^5$ s, when the temperature dropped down to $T_{QCD} \sim \Lambda_{QCD} \sim 300$ MeV. Confinement of colour takes place in this period. The quark-gluon plasma transformed into the gas of colourless hadrons. A detailed picture of this transformation has not been elaborated yet. However, it may be expected that at $T < T_{QCD}$ hadrons were present in the Universe in their usual (observational) form, and not in the form of quarks and gluons. At higher temperatures $T > T_{QCD}$ the density of relativistic quarks and gluons was determined by the equilibrium law, being given by eq. (2.9). After t_{QCD} pions, nucleons and antinucleons were formed, and the baryon excess, generated in the baryosynthesis, was transformed into the excess of baryons over antibaryons. At $t \geqslant 10^{-4}$ s, *i.e.* at $T \leqslant 100$ MeV, pions decay and baryon-antibaryon annihilation proceeds. Due to annihilation practically all $\mathcal{N}\bar{\mathcal{N}}$ pairs were burned and only a small baryon excess $\Delta n_B/n_\gamma \approx 10^{-9\pm1}$ which was created at $t \sim t_{GUT}$ survives.

By the first second of expansion (*i.e* at $T \sim 1$ MeV) approximately equal amounts of photons, neutrinos and antineutrinos of all kinds, electrons and positrons were present in the Universe. There exists a small (about $10^{-8} \div 10^{-10}$ in density) admixture of nucleons. This small admixture is the main prediction of GUT cosmology. It is this very admixture that is the basis of all the visible matter surrounding us.

We may conclude that, in fact, any modification of the presented scenario of the evolution of the very early Universe is possible, provided that it results in the same baryon-to-photon ratio. Thus any other relic of the very early Universe surviving to the successive stages of expansion would be of great importance.

The period $1\ \text{s} \leqslant t \leqslant t_{RD} \sim 10^{12}$ s refers to the radiation dominance (RD) stage of the evolution of the Universe, since the radiation energy density exceeds the matter density in this period. The respective interval of temperature $1\ \text{MeV} \geqslant T \geqslant 1$ eV provides the conditions for the well-established nuclear and atomic processes. Any possible uncertainty in their description is connected with the uncertainty in the magnitude of the respective parameters. The astrophysical impact of these processes leads to observable effects. Thus the principal features of the evolution at the RD stage may be checked in astronomical observations, since the primordial chemical composition and the observed electromagnetic thermal background are relics of this stage.

At $t \sim (0.1 \div 1)$ s (*i.e.* at $T \sim (3 \div 1)$ MeV), the time scale of reactions of weak interaction exceeds the cosmological time scale, so that neutrinos decouple from the other particles and the β reactions

$$\nu_e + n \rightleftarrows e^- + p\,, \qquad \bar{\nu}_e + p \rightleftarrows e^+ + n \tag{2.17}$$

are « switched off ». This means that the ratio of neutron and proton concentrations is frozen out and does not change until neutrons decay at $t \sim 10^3$ s. However, most of the neutrons do not succeed in decaying, since at $t \sim 10^2$ s (at $T_D \sim 100$ keV) they combine with protons forming deuterium in the reaction

$$n + p \to d + \gamma\,. \tag{2.18}$$

At $T > T_D$ the inverse reaction

$$\gamma + d \to n + p \tag{2.19}$$

destroys the bulk of the produced deuterium, but at $T < T_D$ this reaction is inefficient, so the formed deuterium participates in successive thermonuclear transformations:

$$d + d \begin{cases} \nearrow T + p\,, & T + d \to {}^4He + n\,, \\ \searrow {}^3He + n\,, & {}^3He + n \to T + p\,. \end{cases} \tag{2.20}$$

As a result of these transformations the « primordial chemical composition » is formed. This composition is of great importance for our future discussion. Thus a detailed consideration of present observational data on the light-element abundances and the dependences of the calculated primordial abundances of these elements are given in the appendix.

At $T_{rec} \sim 3000$ K ($t_{rec} \sim 10^{13}$ s) recombination of protons and electrons takes place. Neutral atoms are formed and photons decouple from matter. At

$$T_m \sim \frac{1}{3}\frac{n_B}{n_\gamma} m_p \sim \frac{1}{3} r_B m_p \tag{2.21}$$

the density of matter $\varrho_B = m_p n_B$ exceeds the density of photons $\varrho_\gamma \approx n_\gamma 3T$. The moment t_m, when the matter dominance $\varrho_B > \varrho_\gamma$ begins, is determined by the magnitude $r_B = n_B/n_\gamma$ and is rather close to the period of recombination. So at $t > t_{RD}$ (see fig. 2.1) the stage dominated by the modern matter begins. At this stage the gravitational instability of the atomic gas evolves. Small density fluctuations of the matter density grow into the observed structure of inhomogeneities (clusters of galaxies, galaxies, etc.).

The process of growth of density fluctuations takes a rather long time. Long after t_{RD}, matter expands almost uniformly, and the growth of inhomogeneities is reduced to the growth of the density contrast between different regions. Only almost recently at $t \gtrsim 10^{16}$ s the first inhomogeneities, separated from the general cosmological expansion, were formed. Their successive evolution has resulted in the galaxy formation. In the interiors of stars, formed in the galaxies, thermonuclear reactions result in heavy-element production. Stellar explosions at the end of stellar evolution eject heavy elements into the interstellar space. So at the stage of galaxy formation stellar nucleosynthesis takes place.

Based on the physical laws well proved in the laboratories the picture of the cosmological evolution after the first second of expansion presents some definite quantitative predictions.

The general consistency of these predictions with the observations makes the whole picture reliable and seems to leave rather little room for its possible changes. However, recent (see [19, 20]) investigations of the astronomical impact of the neutrino rest mass have shown that big-bang neutrinos—relics from the first second of expansion—can change drastically the whole picture of the successive cosmological evolution. And this change results in even better consistency with the observations.

2'3. *The paradox of antimatter.* – According to the above-presented scenario of the big-bang Universe, the presence of antimatter in the very early Universe (at $t \leqslant t_{\mathrm{QCD}}$) was inevitable. The amount of antiparticles was almost equal to the amount of particles. But in the course of the successive expansion almost all the antiparticles annihilated with the respective particles, so that only a small excess of particles generated by GUT processes survived. The presented scenario leaves practically no room for any sizable amount of antimatter (or antinucleons) at successive stages of the expansion. However, slight modifications of the GUT scenario may result in the prediction of the «late» (*i.e.* long after t_{QCD}) appearance of a considerable amount of antimatter in the Universe. Two possible sources of the « late » appearance of antinucleons in the Universe may be realized: *a*) Survival of antinucleons from the early stages due to suppression of their annihilation in the period after t_{QCD}. (The following possibility exists. If the baryon excess was distributed inhomogeneously, regions with a deficit of baryons may arise, so that antibaryons survive in these regions.) *b*) « Late » production of antinucleons. Decays of frozen heavy metastable particles H of the type

$$\mathrm{H} \to \mathcal{N}\overline{\mathcal{N}} + \text{anything} \tag{2.22}$$

or evaporation of primordial black holes, containing $\mathcal{N}\overline{\mathcal{N}}$ pairs among the products of evaporation, are possible examples of the sources of late $\overline{\mathcal{N}}$ production.

These sources of late appearance of antinucleons in the Universe and their relationship with GUT parameters will be discussed in detail in sect. **6-8**. In any case the appearance of a considerable amount of antinucleons (antimatter) in the Universe long after t_{QCD} may imply substantial astrophysical effects. These effects are considered in the following sections **3** and **4**.

3. – The searches of antimatter.

For the reader's convenience we start by summarizing the present state of the searches for antimatter in the Universe:

1) The standard point of view is that, at present, there is no antimatter on a macroscopic scale up to the scales of clusters of galaxies. This conclusion comes from the observational data of γ-astronomy.

2) The standard cosmological theories lead to the conclusion that, at present, there must be practically no antimatter on a macroscopic scale.

3) Nevertheless, antiprotons in cosmic rays were found and their spectrum is not well consistent with standard models of cosmic-ray propagation. The data of γ-astronomy, in principle, do not contradict the results of calculations based on the assumption of the annihilation nature of the γ-spectrum.

4) Observational restrictions on the amount of annihilated antimatter at early stages of the cosmological evolution are rather weak.

Our point of view is that, very probably, there is no substantial amount of antimatter at present, but this could have possibly been the case in the past (*).

In this section we start from the consideration of results of the direct searches of pieces of antimatter (such as $\overline{\mathrm{p}}$, $\overline{\mathrm{A}}$, antimeteors, etc.), then we discuss the results of indirect searches of antimatter due to the observation of annihilation products. At the end we touch briefly the troubles of the standard baryon-symmetric model of the Universe.

3·1. *Antimatter in the cosmic rays.* – In this subsection we consider the question: did we see pieces of antiworld such as antiparticles, antinuclei or more complex antibodies on the Earth or in its vicinity?

Antiparticles. We have really observed antiprotons in cosmic rays. All observations have been done by a balloon-borne equipment. The results of these experiments are summarized in table 3.I.

(*) To be exact, at this point, we are talking of the presence of considerable amounts of antibaryons in the Universe much later than the so-called «hadronic stage», *i.e.* at $t \gg 10^{-3}$ s.

TABLE 3.I.

Group	$N_{\bar{p}}$ number of $\bar{p}$	$T_{\bar{p}}$ energy of $\bar{p}$ (GeV)	$N_{\bar{p}}/N_p$
1) GOLDEN *et al.* [21]	28	$4.7 \div 11.6$	$(5.2 \pm 1.4) \cdot 10^{-4}$
2) BOGOMOLOV *et al.* [22]	2	$2 \div 5$	$(4 \pm 3) \cdot 10^{-4}$
3) BUFFINGTON *et al.* [23]	14	$0.13 \div 0.32$	$(2.2 \pm 0.6) \cdot 10^{-4}$

These antiprotons did not necessarily come from the antiworld. More probably they have been created by the interaction of cosmic rays with the interstellar medium. The corresponding reaction is well known:

$$p + p \rightarrow \bar{p} + p + p + p + \text{anything}\,.$$

The problem is that the results of observations [23] are not consistent with ordinary cosmic-ray theories in which it is assumed that $\bar{p}$'s are secondaries. In fig. 3.1 we show the results of the theoretical calculations of the $\bar{p}$ flux

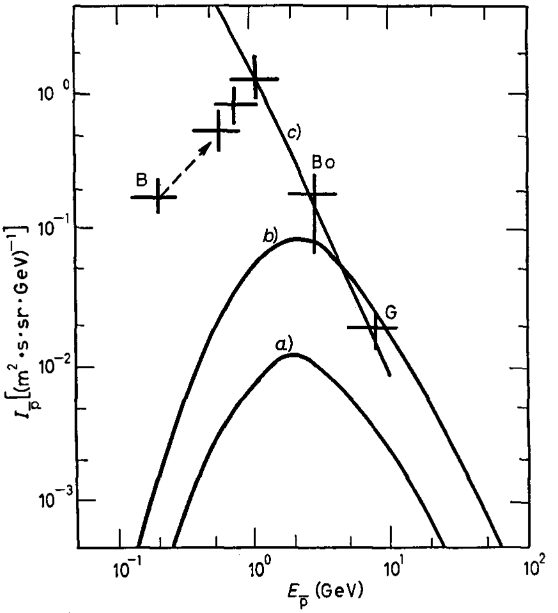

Fig. 3.1. – Secondary-antiproton flux predictions (from [24]). Curve *a*) corresponds to a standard «leaky box» model of cosmic-ray propagation. Curve *b*) corresponds to the same model but the quantity of matter with which the primary protons interact is 16 times bigger. Curve *c*) is a spectrum proportional to that of protons, but arbitrarily scaled down by a factor 2000. Bo denotes Bogomolov *et al.* point [22], G denotes Golden *et al.* [21], B denotes Buffington *et al.* [23]. Three crosses in the upper part correspond to the Buffington point corrected for the effect of solar modulation which tends to reduce the energy of cosmic rays entering the solar system.

from paper [24]. Curve *a*) corresponds to a standard model of cosmic-ray propagation (the so-called « leaky box » model). Curve *b*) is the result of calculations in the framework of the standard model, but under the assumption that the reduced quantity of matter in our Galaxy with which the primary protons interact is 16 times bigger than usually assumed [25]. Curve *c*) is a spectrum proportional to that of protons, but arbitrarily scaled down by a factor 2000. From the results in fig. 3.1 one can conclude that the observational data (or, at least, Buffington's point) are not in satisfactory agreement with the theoretical predictions.

Among the possible explanations of this contradiction, except the trivial one (experiments are wrong or theories are not suitable), we shall mention the hypothesis on n-$\bar{\text{n}}$ oscillations and on the « true » antimatter origin of the observed $\bar{\text{p}}$'s. The n-$\bar{\text{n}}$ oscillations due to baryon number nonconservation, predicted by GUT, induced the transition of the ordinary neutrons to antineutrons which subsequently decayed into $\bar{\text{p}}$'s. However, detailed calculations [26] based on this interesting suggestion seem to rule out this possibility.

STECKER suggested that the cosmic-ray $\bar{\text{p}}$'s are antiprotons from antigalaxies. In a number of papers [9, 27, 28] he favoured the baryon-symmetric cosmology (BSC). According to the BSC matter and antimatter are separated at the level of clusters of galaxies. High-energy particles might be able to leak into neighbouring regions, thus giving rise to an antimatter component of cosmic rays. The shape of the expected $\bar{\text{p}}$ spectrum might not differ from curve *c*) in fig. 3.1. But this would be the case if, due to intergalactic propagation, there were no significant distortions and 1 GeV particles could, finally, reach another cluster of galaxies. One more difficulty of this model arises from the fact that antihelium nuclei $\bar{\alpha}$ are not observed in cosmic rays at present.

Evaporating primordial black holes (PBHs) are another possible exotic source of antiprotons (see sect. **6**). According to HAWKING [29], PBHs with a mass of 10^{15} g are evaporating now, radiating particles as a black body with surface temperature ~ 10 MeV. In the course of evaporation the mass of the PBH decreases, thus increasing the surface temperature of the PBH

$$T_{\text{PBH}} = \frac{10^{13}\ \text{GeV}}{M(\text{g})}\,.$$

So, when less than 1 % of the initial (10^{15} g) mass is left, T_{PBH} exceeds 1 GeV and p$\bar{\text{p}}$ pairs may be radiated by the PBH. PBHs should be distributed in the halo of the Galaxy, so they would be the galactic source of $\bar{\text{p}}$'s. Note that $\bar{\alpha}$ production is strongly suppressed in this mechanism, so there is no difficulty in explaining the absence of $\bar{\alpha}$ in cosmic rays.

Antinuclei. The upper limit of the $\bar{\alpha}/\alpha$ ratio is $2.2\cdot 10^{-5}$ [23]. For other antinuclei the corresponding limits vary from $1\cdot 10^{-2}$ to $9\cdot 10^{-5}$ (see [6, 30]).

No antinucleus has ever been found in cosmic rays, despite the fact that the BS cosmology is not excluded on the basis of these data. It is seen that the observational limits are not very stringent and the detection of even one antinucleus (for example, $\bar{\alpha}$) may turn over all our speculations. We know that the probability to form $\bar{\alpha}$ in proton collisions is extremely small $\sim 10^{-11}$. On the other hand, $\bar{\alpha}$'s may be copiously produced in antistars.

Antibodies. – We put aside the speculative hypothesis of Alfvén [31] that observed γ-bursts are the result of anticomets falling down on stars and mention only the experiments make by KONSTANTINOV *et al.* [32]. They investigated the correlation between the appearance of meteor showers and the increase of the intensity level of γ-radiation and neutrons. They observed an increase of the γ-radiation and of the number of neutrons at altitudes of $(13 \div 18)$ km the moment meteors appear. The effect is about 2 % greater than the background. From the statistical point of view, the effect is greater than the background of 6σ.

3'2. *The observation of products of possible annihilation.* – As well as the proof of the pudding is in eating, the proof of the antiworld is in the annihilation. So, the question is: can the observed data on the possible products of antimatter annihilation result in the conclusion of the existence of antimatter?

3'2.1. Annihilation at rest. Let us consider $p\bar{p}$ annihilation in cosmic space. The typical kinetic energies of the particle we are interested in are small.

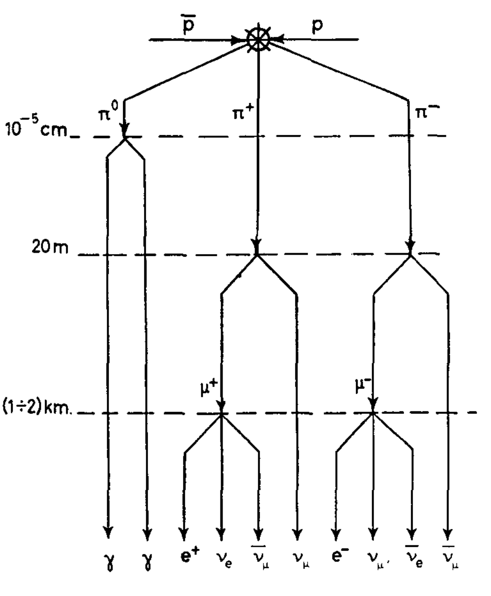

Fig. 3.2.

In table 3.II we summarize the data on multiplicity and branching ratios for different annihilation channels. In figure 3.2 we show schematically the annihilation in space. The numbers on the left-hand side correspond to the mean distance which a particle can travel in space without interaction up to its decay.

TABLE 3.II. – *The main features of the* $p\bar{p}$ *annihilation at rest* [33].

1) Average multiplicity	
$n_{\pi^\pm}$	$3.05^{+0.04}_{-0.036}$
n_{π^0}	1.96 ± 0.23
n_π (total)	5.01 ± 0.23
2) Average energy of pions	
E_π	234 MeV
3) The main annihilation channels	
$\pi^+\pi^-$	0.375 ± 0.03
$\pi^+\pi^-\pi^0$	6.9 ± 0.35
$\pi^+\pi^-X^0$	35.8 ± 0.8
$2\pi^+2\pi^-$	6.9 ± 0.6
$2\pi^+2\pi^-\pi^0$	19.6 ± 0.7
$2\pi^+2\pi^-X^0$	20.8 ± 0.7
$3\pi^+3\pi^-$	2.1 ± 0.25
$3\pi^+3\pi^-\pi^0$	1.85 ± 0.15
$3\pi^+3\pi^-X^0$	0.3 ± 0.1

One can see that at a distance of about 2 km from the point of annihilation only γ-quanta, e^+, e^- and neutrinos survive. An exact calculation shows that a large part of the annihilation energy ($\sim 50\,\%$) is carried away by neutrinos, about $34\,\%$ by γ-quanta and the remaining $16\,\%$ is divided between e^+ and e^-.

So, we, in principle, can search for the signal from annihilation in fluxes of neutrinos, γ-quanta and e^+e^-.

From the beginning we discard from our consideration the e^+e^- component. A lot of e^+e^- pairs is created in the interaction of ordinary cosmic rays with the interstellar medium and with the Earth atmosphere. That background is so large that we cannot extract the e^+e^- from the antimatter annihilation only (for further consideration, see [6]).

3.2.2. Neutrinos from the antiworld. It is well known that at present we have not suitable detectors for extragalactic neutrinos, but this does not exclude the apperance of such detectors in the near future. So, we discuss here the principal possibility to use neutrinos as a signal from the antiworld.

It is easy to show that, in principle, one can distinguish between neutrinos of stars and those of antistars. For example, neutrinos are radiated from stars in the process of deuterium formation

$$\mathrm{p}+\mathrm{p} \to \mathrm{d}+\mathrm{e}^{+}+\nu_{\mathrm{e}}\,,$$

while antistars emit antineutrinos

$$\bar{\mathrm{p}}+\bar{\mathrm{p}} \to \bar{\mathrm{d}}+\mathrm{e}^{-}+\bar{\nu}_{\mathrm{e}}\,.$$

Another important possibility arises due to the collapse of an ordinary star into a neutron star, which is accompanied by huge neutrino radiation owing to the processes of neutronization in the collapsing stellar core:

$$\mathrm{e}^{-}+(A, Z) \to \nu_{\mathrm{e}}+(A, Z-1)\,.$$

The total number of excessive ν_{e} in the course of neutron star formation is easily estimated from the total number of protons in nuclei, which are converted into neutrons via neutronization:

$$N_{\nu_{\mathrm{e}}} \sim N_{\mathrm{p}} \sim \frac{Z}{A}\frac{M_{\odot}}{m_{\mathrm{p}}} \sim 10^{57}\,.$$

In the case of antineutron star formation the same number of excessive $\bar{\nu}_{\mathrm{e}}$ is to be radiated. So, a detailed study of cosmic neutrino fluxes may provide the distinction of neutron and antineutron star formation.

Several suggestions have been made recently for using high-energy neutrinos produced by the interaction of cosmic-ray p (or $\bar{\mathrm{p}}$) with the universal background radiation (see [9, 34]). The typical reactions are

$$\begin{array}{l} \mathrm{p}+\gamma \to \mathrm{n}+\pi^{+} \\ \qquad\qquad\qquad \downarrow \\ \qquad\qquad\qquad \mu^{+}+\nu_{\mu} \\ \qquad\qquad\qquad \downarrow \\ \qquad\qquad\qquad \mathrm{e}^{+}+\nu_{\mathrm{e}}+\bar{\nu}_{\mu}\,, \end{array} \tag{3.1}$$

or

$$\begin{array}{l} \bar{\mathrm{p}}+\gamma \to \bar{\mathrm{n}}+\pi^{-} \\ \qquad\qquad\qquad \downarrow \\ \qquad\qquad\qquad \mu^{-}+\bar{\nu}_{\mu} \\ \qquad\qquad\qquad \downarrow \\ \qquad\qquad\qquad \mathrm{e}^{-}+\bar{\nu}_{\mathrm{e}}+\nu_{\mu}\,. \end{array} \tag{3.2}$$

Therefore, the proton interaction with the relic background radiation results in ν_e formation, while the antiproton interaction leads to $\bar{\nu}_e$. These antineutrinos may interact with the electrons of the Earth atmosphere and for ultrahigh $\bar{\nu}_e$ the cross-section of this interaction can be enhanced due to the formation of weak intermediate bosons W^-:

$$\bar{\nu}_e + e^- \to W^- \to \bar{\nu}_e + e^- .$$

Therefore, the detection of the cosmic $\bar{\nu}_e$ flux may be an indication of the antimatter existence on the macroscopical level. Unfortunately, detailed calculations of the competing background fluxes of $\bar{\nu}_e$ (see [9, 35]) show that the situation is not quite clear, because the background flux of $\bar{\nu}_e$ is comparable with that from the $\bar{p}\gamma$ interaction.

3'2.3. The data of γ-astronomy. The modern γ-astronomy is a very young and rapidly developing branch of astronomy. Many important results have been obtained just within the last few years. Since the Earth's atmosphere is not transparent to γ-rays, almost all observational data were obtained during satellite flights. It was discovered that the greater part of the γ-radiation comes from our Galaxy. But it was shown that the isotropic γ-background exists, too. This flux of γ-radiation comes from all directions and may be generated by processes outside the Galaxy or our cluster of galaxies. We are interested mainly in this γ-background because just from these data one draws the conclusion that macroscopic quantities of antimatter up to the scales of clusters of galaxies do not exist. We discuss briefly how this conclusion was obtained.

The magnitude which is measured is the γ-ray flux N. The absolute value of N is rather small, for $E_\gamma \sim 100$ MeV $N \sim 10^{-5}$ quanta/cm$^2\cdot$s$\cdot$sr. A number of physical processes in cosmic space may induce the γ-radiation: for example, Compton scattering of photons, bremsstrahlung of electrons, interactions of cosmic-ray protons and nuclei with the interstellar medium, etc. Therefore, we cannot determine the nature of γ-quanta, but, if we assume that γ-rays are produced in annihilation, we obtain an upper limit of the possible amount of annihilating antimatter.

It is clear that the γ-ray flux N is proportional to the fraction of antimatter f involved in the annihilation, to the density ϱ of matter in a given space region and to the annihilation rate $\langle\sigma_{\text{ann}} v\rangle$:

$$N \sim f\varrho^2 \langle\sigma_{\text{ann}} v\rangle . \tag{3.3}$$

The annihilation rate is known from experiments, ϱ can be estimated from astronomical observations. In table 3.III the typical estimates for f are summarized.

TABLE 3.III. – *The possible fraction of antimatter f from γ-ray observations* (*according to* [6]).

1) Our Galaxy	$f \leqslant 10^{-15} \div 10^{-10}$
2) Hot intergalactic gas	$f \leqslant 10^{-7}$
3) Clusters of galaxies	$f \leqslant 10^{-5}$
4) Antistars	$f \leqslant 10^{-4}$

One can conclude that antimatter hardly exists in our Galaxy. The limits in the first row of table 3.III are rather stringent. But this conclusion is not exact. If f is small, it may be because the antimatter fraction itself is small or because the fraction of annihilated antimatter is small. So, one may interpret the results of table 3.III as an indication that antimatter and matter are well separated. From this point of view it is interesting to discuss the result in the fourth row. It is obtained from the same measured flux of γ-rays as for the first row, but under the arbitrary assumption that antimatter in some unknown way may form antistars. Then the annihilation region is rather small and the existence of a considerable number of antistars just in our Galaxy is not excluded. ($f \leqslant 10^{-4}$ corresponds to $\sim 10^7$ antistars (*).)

As we mentioned above, the data in table 3.III are obtained from the observation of the background γ-ray radiation. If the value of f for clusters of galaxies is $f > 10^{-5}$, then such clusters would be discrete observable γ-ray sources, but not the background. Up to now we have not been able to detect such extragalactic sources of γ-radiation. So, one concludes that, if antimatter exists, it is separated from matter on the scales of clusters of galaxies.

An interesting result has been obtained by STECKER [9, 27]. In fig. 3.3 the energy spectrum of the background radiation is shown. At a first glance this spectrum is quite inconsistent with the expected energy distribution for γ-quanta from decays of annihilation π^0's.

Nevertheless, STECKER has shown that, if the red-shift of photon energy induced by the expansion of the Universe as well as the absorption of photons in the interstellar gas are taken into account, the shape of the γ spectrum from π^0 decay must change substantially.

The most important free parameter in the calculations is the time (or red-shift Z) at which the annihilation takes place. The results of the calculations coincide with the observational data, when $Z \sim 100$. This corresponds to the annihilation which has proceeded to the time of $t \sim 10^7$ y from the beginning

(*) We would like to stress that this example in no way can be considered as a « proof » of antistar existence. The only thing we intend to show is that one must keep in mind that the possibility of sharp boundaries between matter and antimatter results in a decrease of the value of f.

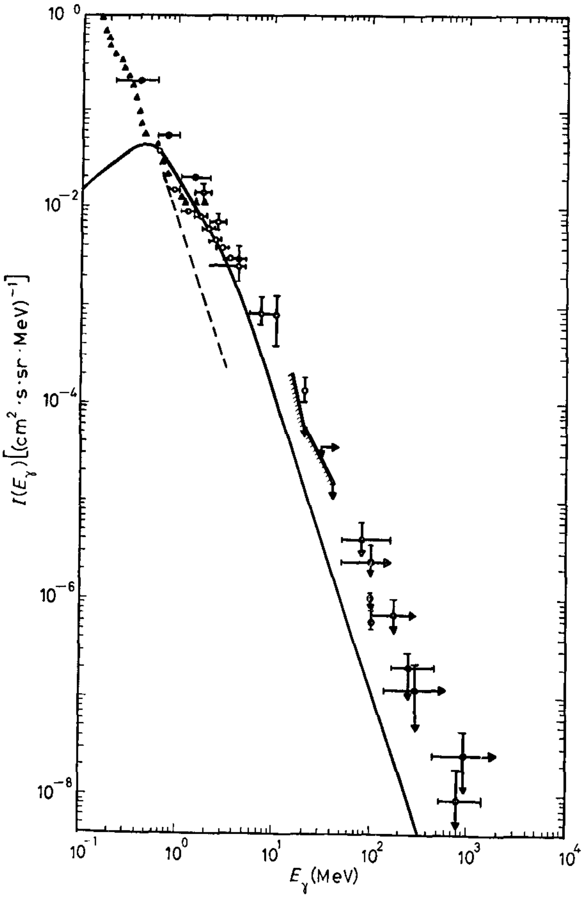

Fig. 3.3. – Cosmic γ-ray background spectrum from matter-antimatter annihilation and observational data (from [27]).

of the expansion. The full line in fig. 3.3 represents the results of those calculations. It is seen that the assumption about the annihilation nature of the γ-ray background spectrum does not contradict the observational data.

3'3. *Annihilation at the stage of radiation dominance in the Universe.* – Up to now we have been discussing mainly the possibility of annihilation at the present time. It is quite probable that now antimatter does not exist in macroscopical amounts, but antimatter may have existed in the early Universe. The presence of antibaryons in equilibrium at $t < 10^{-6}$ s ($T > 1$ GeV) was inevitable. The question is: Could a sizable amount of antibaryons manage to survive up to later periods? What restrictions can be put on late annihilation?

Modern developments of the theories of elementary particles and astrophysics provide a number of mechanisms of conservation or late production of antibaryons. These sources of antimatter could be evaporation of primordial black holes, decays of heavy particles, domains of antimatter, etc. We will discuss the properties of these sources later, in sect. **6-8**.

In this subsection we stand on a purely phenomenological point of view and put aside the question on how antimatter appeared in the Universe. We shall consider the possibility that antimatter has not survived up to the present time because of its annihilation with matter. Evidently, such a possibility may arise within the frame of the baryon-asymmetric model only.

So, how may we check the presence of antimatter at the earlier stages of the cosmological expansion?

The first possibility, again, is in close touch with γ-rays from annihilation. Being red-shifted due to the expansion, they survive until the present time. But at red-shifts $Z > 10^2$ the Universe was opaque to γ-rays, so that the energy released in the annihilation heats the matter. Hot electrons appear in the Universe and their interaction with the electromagnetic background induces distortions of the spectrum of the relic radiation. The quantitative theory of the distortions of the relic-radiation spectrum is given in [1, 6]. We shall give here only some brief discussion of their results concerning the limits on the possible amount of annihilated antimatter.

The theory (cf. [1]) of distortions of the relic-radiation spectrum considers two different cases (depending on Z): *a*) early energy release at

$$10^8\,\Omega_b^{\frac{1}{2}} > Z > 4\cdot 10^4\,\Omega_b^{\frac{1}{2}} \tag{3.4}$$

and *b*) late energy release at

$$4\cdot 10^4\,\Omega_b^{\frac{1}{2}} > Z\,. \tag{3.5}$$

If the energy was released at $Z > 10^8\,\Omega_b^{\frac{1}{2}}$ (*i.e.* at $t < 10^3$ s), the reactions

$$e+p \to e+p+\gamma\,, \qquad \gamma+e \to 2\gamma+e \tag{3.6}$$

could provide the formation of an additional number of photons, so that the Planck equilibrium spectrum of photons succeeds in setting up after the energy has been released. This means that any energy release at $Z > 10^8\,\Omega_b^{\frac{1}{2}}$ does not induce distortions of the thermal spectrum. But at $Z < 10^8\,\Omega_b^{\frac{1}{2}}$, reaction (3.6) is ineffective in producing additional photons, its effect is negligible. Thermal equilibrium between hot electrons and cold photons sets up under the condition of a constant number of photons. So in case *a*), *i.e.* if the energy has been released in period (3.4), the photon distribution, maintained after equilibrium between photons and electrons is established, is not the thermal Planck distribution (2.3), but is the thermal distribution for a fixed total number of photons, *i.e.* the Bose-Einstein distribution

$$F_\nu = \frac{2h\nu^3}{c^2} \frac{1}{\exp[(h\nu + \mu)/kT_e] - 1}, \tag{3.7}$$

where ν is the frequency, T_e is the electron temperature and the chemical potential μ is determined by the relative magnitude of the energy release. The related deviation of the formed spectrum (3.7) from the Planckian distribution induces distortion of the thermal-background spectrum. The upper limit on the possible energy release from the observed spectrum of relic radiation is in the considered case

$$\left(\frac{\delta\varepsilon}{\varepsilon_\gamma}\right)_a < 6 \cdot 10^{-2}, \tag{3.8}$$

where $\delta\varepsilon$ is the specific energy release and ε_γ is the photon energy density.

However, if the energy release is late (*i.e.* at $Z < 4 \cdot 10^4 \Omega_b^{\frac{1}{2}}$), the equilibrium Bose-Einstein spectrum (3.7) does not succeed in setting up. The distortions of the photon spectrum are determined by the kinetics of the heating of the photon gas by hot electrons. The observational upper limit on such distortions implies the following upper limit on the late energy release inducing such distortions

$$\left(\frac{\delta\varepsilon}{\varepsilon_\gamma}\right)_b < 2 \cdot 10^{-1}, \tag{3.9}$$

where $\delta\varepsilon$ is the specific late energy release, ε_γ is the energy density of photons and index the b signifies that case (3.5) is considered.

Based on restrictions (3.8) and (3.9) on the relative energy release we may obtain restrictions on the relative amount f of annihilated antinucleons. Indeed, the photon energy density ε_γ is given by

$$\varepsilon_\gamma = 6 \cdot 10^{-13} \left(\frac{T_0}{3\text{K}}\right)^4 (1 + Z)^4 \ \text{erg/cm}^3, \tag{3.10}$$

where T_0 is the temperature of relic photons at the present time. The energy release $\delta\varepsilon$ is determined by the relative amount f of annihilated antinucleons and is given by

$$\delta\varepsilon = 5 \cdot 10^{-9} \Omega_b f (1 + Z)^3 \ \text{erg/cm}^3, \tag{3.11}$$

where $\Omega_b = \varrho_b/\varrho_c$ is the baryon density (in units of the critical density). From

(3.10) and (3.11) one obtains

$$\frac{\delta\varepsilon}{\varepsilon_\gamma} = \frac{8\cdot 10^3 \Omega_b f}{1+Z}\left(\frac{3\mathrm{K}}{T_0}\right)^4 . \tag{3.12}$$

So upper limits on the distortions of the thermal-background spectrum result in the following upper limits on the value of f:

$$f \leqslant \left(\frac{T_0}{3\mathrm{K}}\right)^4 \frac{1}{8\cdot 10^3 \Omega_b}(1+Z)\begin{cases} 2\cdot 10^{-1} & \text{at } Z < 4\cdot 10^4 \Omega_b^{\frac{1}{2}} , \\ 6\cdot 10^{-2} & \text{at } 4\cdot 10^4 \Omega_b^{\frac{1}{2}} < Z \leqslant 10^8 \Omega_b^{\frac{1}{2}} . \end{cases} \tag{3.13}$$

We see that the limits of f are quite weak at $Z \geqslant 10^3$ (*i.e.* at $t < 10^{13}$ s):

$$f \leqslant 1 \quad \text{at } \Omega_b \sim 0.1 .$$

4. – Annihilation with ^{4}He. The best test of existence of antimatter in the early Universe.

Let us consider the period after the end of the primordial big-bang nucleosynthesis but before the recombination of hydrogen, *i.e.* $10^3\,\mathrm{s} \leqslant t \leqslant 10^{13}$ s. As we have discussed in subsect. **3**.3, the limits on the possible amount of antimatter at this stage of the cosmological evolution obtained from the distortion of the thermal relic background are quite weak, $f < 1$. It is easy to obtain a more stringent restriction on f from the investigation of the antiproton annihilation with ^{4}He.

As we mentioned in sect. **2**, ^{4}He is the most abundant element in the Universe after hydrogen. Its concentration in weight $X_{^4\mathrm{He}} = n_{^4\mathrm{He}}/n_{\mathrm{H}} = 0.24$, while the concentrations in weight of the other elements are considerably small. For example, $X_{\mathrm{D}} \simeq 2.5\cdot 10^{-5}$ and $X_{^3\mathrm{He}} = 4.2\cdot 10^{-5}$ [36]. When antiprotons annihilate with ^{4}He (just at rest), they may create deuterium and ^{3}He in the reactions

$$\bar{\mathrm{p}} + {}^4\mathrm{He} \to {}^3\mathrm{He} + N(\pi) , \; \mathrm{d} + \mathrm{p(n)} + N(\pi) , \tag{4.1}$$

where N is the number of pions. Therefore, if antimatter did really exist in the early Universe after the big-bang nucleosynthesis, this would inevitably lead to the formation of D and ^{3}He. From the comparison of the concentrations in weight of ^{4}He, D and ^{3}He it is easy to see that the destruction of a quite small part of ^{4}He ($\sim 10^{-4}$) in annihilation may create all the observed abundance of D and/or ^{3}He.

It must be noted that deuterium may be formed due to annihilation in the early Universe not only in the direct reactions of the type of (4.1). As we mentioned in sect. **2**, there were no free neutrons remaining in the Universe after the end of the primordial nucleosynthesis (at $t \geqslant 10^3$ s): all neutrons were either hidden in nuclei or decayed. Due to $\bar{p}^4$He annihilation a number of free neutrons appear in the reactions

$$\bar{p}+{}^4\text{He} \to N_1(\text{n})+N_2(\text{p})+N_3(\pi)\,, \tag{4.2}$$

where N_1, N_2 and N_3 are the numbers of neutrons, protons and pions.

If the proton density in the Universe at the moment of annihilation is sufficiently high, then the neutrons may succeed in colliding with the protons before decay, forming deuterium in the reaction

$$\text{n}+\text{p} \to \text{d}+\gamma\,. \tag{4.3}$$

Therefore, an additional amount of deuterium is created.

So, there are two processes of the deuterium formation due to $\bar{p}^4$He annihilation: 1) the direct process and 2) the indirect one. The former takes place any time when annihilation occurs, but the latter may proceed only in the hot early Universe. The indirect mechanism dominates, when it is possible. That is simply because the number of neutrons in $\bar{p}^4$He annihilation is greater than the number of deuterons.

The indirect mechanism is permitted when the neutrons are mainly captured by protons rather than decay, *i.e.* the time of reaction (4.3) τ is less than the neutron lifetime $t_\text{n} \sim 10^3$ s:

$$\tau = \frac{1}{n_\text{p}\langle\sigma v\rangle} \ll t_\text{n}\,, \tag{4.4}$$

where n_p is the concentration of protons, σ is the cross-section of reaction (4.3) and v is the velocity of n's. The estimate from [37] shows that, in the periods after $t_\text{D} = 4.5\cdot 10^6\,\Omega_\text{b}^{\frac{2}{3}}$ s, the indirect mechanism is suppressed ($\tau > t_\text{n}$). Therefore, when the Universe is older than $t_\text{D} \approx 0.97\cdot 10^6$ s ≈ 3 months (for $\Omega_\text{b} = 0.1$), the only mechanism of deuterium formation is the direct one.

Now let us estimate the limits on f which one may obtain from the study of $\bar{p}^4$He annihilation.

The additional amount of D being created due to $\bar{p}^4$He annihilation is

$$\Delta n_\text{D} = \begin{cases} n_{^4\text{He}}(f_\text{D}+f_\text{n})f & \text{at } 10^3\,\text{s} \ll t \ll t_\text{D}\,, \\ n_{^4\text{He}} f_\text{D} f & \text{at } t > t_\text{D}\,, \end{cases} \tag{4.5}$$

where $n_{^4\text{He}}$ is the concentration of ^{4}He, f_n, f_D are the mean numbers of n and D created in the annihilation. For ^{3}He it will be, correspondingly,

$$\Delta n_{^3\text{He}} = n_{^4\text{He}} f_{^3\text{He}} f\,. \tag{4.6}$$

If we assume that, for example, Δn_D does not exceed the observed abundance of D, *i.e.* X_D, the following restriction on the value f is obtained:

$$f \leqslant \begin{cases} \dfrac{2X_D}{X_{^4He}(f_n + f_D)} & \text{at } 10^3\,\text{s} \leqslant t \leqslant t_D\,, \\ \dfrac{2X_D}{X_{^4He} f_D} & \text{at } t > t_D\,, \end{cases} \tag{4.7}$$

where $X_{^4He}$ is the observed abundance of ^{4}He.

If, as a rough estimate, one assumes that all annihilation channels for ^{4}He have equal cross-sections and $\sigma_{ann} = 0.5\sigma_{tot}$ (where σ_{tot} is the total cross-section of the $\bar{p}\,^4$He interaction), then it is easy to show that

$$f \leqslant 10^{-4}\,. \tag{4.8}$$

Therefore, the experimental investigation of $\bar{p}\,^4$He annihilation may provide valuable information on the outputs of D, ^{3}He and n, so to obtain a limit on the possible amount of antimatter in the early Universe which is at least by three orders of magnitude more stringent than that which comes from distortions of the background spectrum. We shall discuss the experimental situation and some other important astrophysical aspects of $\bar{p}\,^4$He annihilation in sect. **9**.

5. – Antimatter in the baryon-symmetrical Universe.

This section is the first one where we start to discuss the modern theoretical views on the presence of antimatter in the Universe. We shall consider a number of different aspects of this problem. For the reader's convenience we put the final conclusions just here. They are the following:

1) In the framework of a standard (*) baryon-symmetrical cosmology there is no universally acknowledged possibility for the existence of a considerable quantity of antimatter in the Universe at $t > 10^{-3}$ s. Standard baryon-asymmetrical models cannot explain the baryon asymmetry of the Universe, but simply take this fact as an initial condition.

2) The modern cosmology does explain the baryon asymmetry of the Universe. The recent progress in the development of the grand-unification theories (GUT) of elementary particles provides the mechanism for the generation of the baryon excess in the Universe.

3) There are possibilities for the existence of antimatter in the early Universe. They appear in the framework of GUT and in the physics of black holes.

(*) We use here and below the term « standard » for designating the models in the pre-GUT time.

In this section we consider point 1)—the troubles of the baryon-symmetric cosmology.

It is quite natural to suppose that, as every particle has its antiparticle, the same symmetry must occur for the macroscopic parts of the Universe. Let us consider the consequences of this hypothesis for the big-bang cosmology.

5'1. *Antinucleons in the hot primordial plasma (homogeneous baryon-symmetric model).* – As was shown in sect. **2**, the early Universe consisted of a hot plasma where the particles (or matter) were in thermal equilibrium with the radiation (*).

When the temperature of the Universe was higher than $T\sim 1$ GeV, the number of antinucleon-nucleon pairs was as large as the number of γ-quanta, *i.e.* 10^8 times greater than the present amount of protons. This situation occurred at the times $t\leqslant 10^{-6}$ s from the beginning of the expansion of the Universe. The concentration of nucleons and antinucleons depends at this stage on the temperature as follows:

$$n_{\mathcal{N}} = n_{\overline{\mathcal{N}}} = \beta T^3 , \tag{5.1}$$

where β is the number of independent kinds of particles. (We will be working with units where $\hbar = c = k = 1$.)

When the temperature of the Universe was $T < M$, where M is the nucleon mass, the concentration of $\mathcal{N}$ and $\overline{\mathcal{N}}$ decreased exponentially:

$$n_{\mathcal{N}} = n_{\overline{\mathcal{N}}} \approx (MT)^{\frac{3}{2}} \exp[-M/T] . \tag{5.2}$$

It is important to note that (5.2) is valid only for the systems in thermal equilibrium. In the case of $\mathcal{N}\overline{\mathcal{N}}$ pairs this equilibrium occurs due to the balance between $\mathcal{N}\overline{\mathcal{N}}$ annihilation and creation of $\mathcal{N}\overline{\mathcal{N}}$ pairs by γ-quanta. As the annihilation is a quite « strong » reaction and proceeds very fast, the period during which $\mathcal{N}\overline{\mathcal{N}}$ pairs are in equilibrium is quite long.

Exact calculations [1] show that $\mathcal{N}\overline{\mathcal{N}}$ pairs were out of equilibrium only at $t\sim 10^{-3}$ s, when $T\sim 20$ MeV. That is why the stage of exponential decrease is so long, from 10^{-6} s to 10^{-3} s, the concentrations of $\mathcal{N}$ and $\overline{\mathcal{N}}$ fell drastically

$$n_{\mathcal{N}} = n_{\overline{\mathcal{N}}} \approx (10^{-17}\div 10^{-18})\, n_{\gamma} .$$

(*) The picture depicted below is not quite exact, but discussed only from the pedagogical point of view. In the exact considerations of the situation in the Universe when the temperature is higher than 300 MeV one must take into account the dynamics not of the $\mathcal{N}\overline{\mathcal{N}}$ plasma, but of the quark-antiquark one. A consideration of the phase transition from quark plasma to the hadronic stage is needed, too. But the main ideas of this section remain unchanged.

When $t > 10^{-3}$ s these concentrations were « frozen » and practically did not change up to the present time. Therefore, the baryon-symmetrical models (where $n_{\mathcal{N}} = n_{\overline{\mathcal{N}}}$) predict extremly small concentrations of antinucleons as well as of nucleons. The present amount of nucleons is $n_{\mathcal{N}} \approx 10^{-(8\div 10)} n_{\gamma}$, so the « natural » assumption about equal quantities of $\mathcal{N}$ and $\overline{\mathcal{N}}$ in the early Universe leads to disagreement with observation up to ten orders of magnitude.

5·2. *Homogeneous baryon-asymmetrical model.* – As we have shown in the previous subsection, the consideration of the fate of the equal number of $\mathcal{N}$ and $\overline{\mathcal{N}}$ in the early Universe leads to an unsatisfactory conclusion. But we see protons everywhere, whereas antiprotons are observed under very special conditions. So, one can assume that from the beginning there was some excess of baryons over antibaryons. In the first instants of the expansion (10^{-6} s $\leqslant t \leqslant 10^{-3}$ s) all antibaryons annihilated and the excess of baryons formed the whole observed world. As we have seen in subsect. **5·1**, this excess is quite small: it must be an extra proton for 10^8 proton-antiproton pairs. In the standard model the question about the nature of this excess is not considered. It is assumed that the baryon charge of the Universe is nonzero and that is the initial condition.

In recent times an understanding of this problem has been obtained and we will discuss the possible mechanism of generation of the baryon excess in detail in sect. **7**. Now we briefly consider the other possibilities to avoid the troubles of the standard baryon-symmetric model.

5·3. *The unconventional approaches.* – It is clear that, to avoid the intensive annihilation in the early Universe, one must suggest some mechanism for dividing matter and antimatter apart. In principle it is possible to obtain that, if matter and antimatter can, in some way, gather to the regions with different baryon charge, *i.e.* domains of matter and domains of antimatter are formed. Obviously this separation must occur at the very early stage of the Universe expansion $t < 10^{-6}$ s.

The most developed theory of this type is the model of Omnès [38]. Omnès assumes that the hot baryon-antibaryon plasma is not stable and at $t \leqslant 10^{-6}$ s in the early Universe a phase transition did occur that led to the separation of matter from antimatter. Small drops $\sim 10^{-3}$ cm of antimatter and matter were formed. When the temperature of the Universe T decreased, the annihilation between those drops began, but that annihilation was not so strong as for a uniformly mixed $\mathcal{N}\overline{\mathcal{N}}$ plasma, because it proceeded only on the boundaries of the drops. Due to this effect it is possible, in principle, to obtain the correct value for the concentration of baryons $n_{\mathrm{B}}/n_{\gamma} \sim 10^{-8}$. When enough baryons and antibaryons were burned to obtain $n_{\mathrm{B}}/n_{\gamma} \sim 10^{-8}$, the new processes began. According to Omnès, the drops started coalescing one on another. This coalescence proceeded continuously up to the moment of the recombination of hydrogen, *i.e.* $t \sim 10^{13}$ s. The scales of the drops grew enormously up to $\sim 10^{22}$ cm, which is equivalent to the corresponding mass of the domains $\sim 10^{11} M_{\odot}$, *i.e.* as large as for big galaxies.

So, according to OMNÈS, our Universe is, on the average, baryon symmetrical, though large regions of matter and antimatter exist.

There is a great number of deficiencies in the Omnès theory. A full account of these shortcomings one may find in [1, 6]. We list here only few typical ones without any detailed consideration:

1) The existence of the phase transition in the hot baryon-antibaryon plasma at $T \sim 1$ GeV is not proved.

2) Due to the annihilation at $t \sim 10^{13}$ s an appreciable quantity of radiation energy must appear, thus leading to the distortion of the thermal-background spectrum.

3) There are difficulties in explaining the ^{4}He abundance, and the density of galaxies, according to Omnès, is too high as compared to the observed one.

In spite of these troubles, the Omnès idea itself is quite interesting. He makes an attempt to explain the structure of the Universe on the basis of the instability of the hot hadron plasma and its subsequent dynamics without any additional *ad hoc* assumptions. That bold idea explaining the development of the Universe from first principles of the behaviour of elementary particles is quite fruitful. In the next sections we will discuss how this principle works on the basis of modern theories of elementary particles.

6. – Antimatter and inhomogeneities.

At the end of the preceding section we have established the connection between antimatter survival and inhomogeneities of the matter-antimatter distribution. Let us consider this relationship in another aspect: in its connection with the general problem of primordial inhomogeneities, *i.e.* with the problem of initial conditions for the theories of galaxy formation.

In the expanding Universe there had been no observed inhomogeneities at the early stages of the expansion. There had been no stars, no galaxies and their clusters. The homogeneity observed at present in the average distribution of matter, extrapolated into the past, must have transformed into an almost complete homogeneity of the plasma in equilibrium with the radiation. Small density perturbations, present then in the early Universe, must have grown up to the observed structure of inhomogeneities. This is the main point of the modern theory of galaxy formation. The nature of these inhomogeneities may be different. Two approaches to this problem exist. In the first one, for a given density perturbation of matter, related perturbations of radiation density exist, so that the specific entropy of matter (number of photons per nucleon) is not perturbed. These perturbations are called adiabatic [1], and the theory of evolution of such perturbations is called the adiabatic theory (A) of galaxy formation. The second approach considers density perturbations of matter, *i.e.* of the baryon charge, so that the total density of radiation and baryons is constant. In this approach the specific entropy of perturbed and unperturbed regions is different. These perturbations are thus called entropy ones. The theory of evolution of such inhomogeneities is called the entropy (E) theory of galaxy formation [1].

One of the most important questions of both approaches is the question on the initial spectrum of inhomogeneties. Under the notion of spectrum of inhomogeneities we mean the correlation between the amplitude of the perturbation and the size of the perturbed region. The amplitude of the perturbation is defined by the relative magnitude of the deviation of the density in the perturbed region from the average one. The spectrum of inhomogeneities determines the parameters of the structure of inhomogeneities and its evolution, *i.e.* the present and the future picture of the night sky. Some attempts to connect this spectrum with quantum fluctuations in the very early Universe were made recently [39]. Whatever the success of these attempts would be, it is of importance to have observational restrictions on the initial spectrum.

Astronomical observations provide information on the long-wave part of the spectrum, corresponding to the scales of stellar clusters, galaxies, galaxy clusters, superclusters ... up to the modern cosmological horizon. But for a complete picture of the cosmological evolution of the Universe perturbations on all the scales are of importance. From the viewpoint of the statistical theory, the spectrum of fluctuations characterizes the dispersion of the Gaussian distribution of the amplitude. It is the magnitude

$$\sqrt{\overline{(\delta\varrho/\varrho)^2}} = \delta(M) \tag{6.1}$$

that determines the spectrum. Even for small $\delta(M)$ statistical fluctuations provide an exponentially small probability for the existence of density perturbations with amplitude of order 1.

It turns out that fluctuations with amplitude of order 1 in both theories are closely connected with the sources of antimatter. Indeed, the existence of an adiabatic fluctuation with amplitude of order 1 at a certain small scale means that, at a very early stage of the expansion, the region in which such fluctuation arise separates from the cosmological expansion, forming a black hole.

These black holes, formed in the early Universe long before star and galaxy formation, are called primordial in contrast to those black holes which are awaited to be formed in the the final stage of the evolution of massive stars. The possibility of primordial-black-hole (PBH) formation was first discussed by ZEL'DOVICH and NOVIKOV [40] (see also [41]).

As was pointed out by HAWKING [29], PBHs of small mass evaporate (see subsect. **6**.1.1). The energy of particles produced at evaporation is in inverse relationship with the mass of PBHs. Thus small-mass PBHs evaporate relativistic particles, and among the products of evaporation $\bar{p}p$ pairs are to be present. So, owing to the possibility of PBH formation and successive evaporation, small-scale adiabatic fluctuations may be reflected in $\bar{p}$ production.

In the case of entropy fluctuation, a baryon charge fluctuation of order 1 means that at this scale the baryon charge density changes sign, *i.e.* at this scale domains of antimatter arise.

Antimatter domains of sufficient scale may survive to the stage after cosmological nucleosynthesis, thus providing late $\bar{p}$ annihilation.

Mechanisms of PBH and antimatter domain formation in the framework of GUTs will be discussed in sect. **8**. Here we shall discuss observational effects of the existence of PBH and antimatter domains.

6.1. *Evaporating* PBHs *as source of antimatter.* – The most important point for our discussion is the possibility of PBH evaporation discovered by HAWKING [29].

6.1.1 PBH evaporation. HAWKING [29] considered quantum effects in the vicinity of a PBH. A black hole with mass M has radius

$$r_g = \frac{2GM}{c^2} \quad \text{or} \quad r_g = 2GM = \frac{2M}{m_{\rm Pl}^2} \tag{6.2}$$

in units $\hbar = c = 1$.

In the gravitational field of the black hole particles may be created. Particles with energies $E \leqslant 1/r_g$ have wave-length $\lambda \geqslant r_g$, so they may be found beyond the gravitational radius of the black hole, they may escape from the black hole and go to infinity. Thus, owing to quantum effects, radiation from the surface of a black hole is possible. The black hole—the object from which nothing can escape in the classical limit—may radiate. Black-hole radiation is described as thermal black-body surface radiation with temperature

$$T = \frac{1}{4\pi r_g} = \frac{1}{8\pi GM} = \frac{m_{\rm Pl}^2}{8\pi M} = \frac{10^{13}\,\text{GeV}}{M/1\text{g}} \tag{6.3}$$

(units $\hbar = c = k = 1$).

Thus its luminosity is of the order of magnitude

$$\frac{\mathrm{d}\mathscr{E}}{\mathrm{d}t} = \sigma T^4 4\pi r_g^2 \sim \frac{1}{r_g^4} r_g^2 \sim \frac{1}{r_g^2} \sim \frac{m_{\rm Pl}^4}{M^2}. \tag{6.4}$$

The energy release (6.4) implies loss of mass of the black hole ($\hbar = c = 1$)

$$\frac{\mathrm{d}M}{\mathrm{d}t} = -\frac{\mathrm{d}\mathscr{E}}{\mathrm{d}t} = -\frac{m_{\rm Pl}^4}{M^2} \tag{6.5}$$

and during the time interval $t \sim (M/m_{\rm Pl})^3 t_{\rm Pl}$ a black hole of mass M loses all its mass—the black hole evaporates!

With due account of all the numerical factors the time of evaporation is

$$t_e = 10^{-27}\,{\rm s}\left(\frac{M}{1{\rm g}}\right)^3 . \tag{6.6}$$

For black holes of $\geqslant 10 M_\odot$, which are awaited to be formed as a result of stellar evolution, the time scale of evaporation is 10^{74} s $\simeq 10^{66}$ y, so far this massive-black-hole evaporation is negligible. However, the masses of the primordial black holes may be much smaller than the stellar ones. All the values down to the Planck mass $m_{\rm Pl}$ (and even smaller [42]) are possible. For PBHs with mass less than 10^{15} g the time scale of evaporation is smaller than 10^{18} s. Being formed in the early Universe, such black holes must have totally disappeared at the present time. However, the effects of their evaporation may lead to observable consequences, thus providing a definite check of their existence in the past. A detailed discussion of PBHs and restrictions on their concentration in the early Universe may be found in [43]. We shall confine ourselves to considering the antinucleon output of evaporating PBHs.

6.1.2. Antiproton fluxes from evaporating PBHs. Let us estimate the fraction $F_{\bar{p}}$ of total energy evaporated by PBHs in the form of antiprotons. It is seen from eq. (6.3) that evaporating PBHs with mass $M < M_{\bar{p}} = 10^{13}$ g have a surface temperature exceeding 1 GeV, so antiproton production is possible. According to the modern views on hadron production at high energies quarks and gluons are produced first, and afterwards their fragmentation into hadrons takes place. The same seems to be true for PBH evaporation—in the vicinity of an evaporating black hole thermal radiation consists of equilibrium fractions of gluons and quarks. However, at distances of the order $1/\Lambda_{\rm QCD}$ confinement of the colour takes place and quarks and gluons form hadrons. In this picture the antiproton yield is determined by the fragmentation functions of quarks and gluons into $\bar{p}$'s. Minimal estimates of this yield $f_{\bar{p}}$ may be based on considerations [44] of $\bar{p}$ production in e^+e^- annihilation: $F_{\bar{p}} = f_{\bar{p}} \langle n_{\bar{p}} \rangle$ ($\langle n_{\bar{p}} \rangle \sim 0.03$), where $\langle n_{\bar{p}} \rangle$ is the mean $\bar{p}$ multiplicity in e^+e^- annihilation. However, PBH evaporation differs from e^+e^- annihilation by the possibility of direct (equilibrium) production of hard gluons, whose fragmentation onto $\bar{p}$'s is higher than the fragmentation of hard antiquarks and bremsstrahlung gluons produced in e^+e^- annihilation. So, at the evaporation of PBHs with mass $\leqslant M_{\bar{p}} = 10^{13}$ g an $f_{\bar{p}}$ of about $0.05 \div 0.01$ is awaited.

Since evaporation reduces the mass of PBHs, antiprotons may be produced at late stages of evaporation of PBHs with mass $M > M_{\bar{p}} = 10^{13}$ g when their mass has decreased down to $M_{\bar{p}}$. In this case the fraction of total energy carried away by antiprotons amounts to $F_{\bar{p}} = f_{\bar{p}}(M_{\bar{p}}/M)$.

Relativistic antiprotons, evaporating at $t < 10^2$ s (at $T > 100$ keV), are slowed down before annihilation, since, owing to the presence of e^+e^- pairs in equilibrium, the rate of antiproton energy loss is greater than the rate of $\bar{p}p$ annihilation. In the periods after positron annihilation (at $t > 10^2$ s) the density of electrons is small (of the order of the baryon density) and energy losses due to Coulomb interactions with the plasma are small. However, in $\bar{p}p$ interactions at high energies the annihilation channel is suppressed (its magnitude is of the order of the difference between $\bar{p}p$ and pp cross-sections). Reactions of the type $\bar{p}p \to p + \text{anything}$ dominate, decreasing the energy of $\bar{p}$'s. So even at $t > 10^2$ s a substantial part of relativistic $\bar{p}$'s from evaporating PBHs is slowed down. At $t < \Omega_B^2 \cdot 10^{16}$ s the rate of annihilation exceeds the expansion rate. So most $\bar{p}$'s generated in this period succeed to annihilate.

Let us estimate now the density of $\bar{p}$'s produced at $t_e = 10^{-27}\,\text{s}(M/1\text{g})^3$ by evaporating PBHs of mass M. The total amount $N_{\bar{p}}$ of $\bar{p}$'s generated by one such PBH is given by

$$N_{\bar{p}} = \int\limits_0^{\min(M, M_{\bar{p}})} \frac{dN_{\bar{p}}}{dM}\, dM\,, \tag{6.7}$$

where

$$\frac{dN_{\bar{p}}}{dM} = \frac{M}{m_{Pl}^2} f_{\bar{p}}(M)\,. \tag{6.8}$$

The density of $\bar{p}$'s is then given by

$$n_{\bar{p}} = n_{PBH} N_{\bar{p}}\,, \tag{6.9}$$

where $n_{PBH}(M)$ is the concentration of PBHs of mass M in the Universe at the moment t_e of their evaporation:

$$n_{PBH}(M) = \frac{\varrho_{PBH}(M)}{M} = \frac{\alpha(M)\varrho}{M}\,. \tag{6.10}$$

Here $\alpha(M) = \varrho_{PBH}/\varrho$ is the relative contribution to the cosmological density ϱ of PBHs of mass M at the moment of their evaporation. From (6.7)-(6.10) we obtain (assuming $f_{\bar{p}}(M < M_{\bar{p}}) = \text{const}$, $f_{\bar{p}}(M > M_{\bar{p}}) = 0$)

$$n_{\bar{p}} = \frac{\min(M, M_{\bar{p}})}{m_{Pl}^2} f_{\bar{p}} \alpha(M) \varrho\,. \tag{6.11}$$

At the RD stage the mass of evaporating PBHs $M < M_{\bar{p}} = 10^{13}$ g, $\varrho \approx 3Tn_\gamma$ and (see (2.12))

$$T = \left(\frac{45}{32\pi^3}\right)^{\frac{1}{4}} \frac{1}{\varkappa^{\frac{1}{4}}} \sqrt{\frac{m_{Pl}}{t_e}}\,.$$

With the use of (6.5) we have $T \sim \sqrt{m_{Pl}^5/M^3}$. Substituting the expression for ϱ into (6.11), one obtains

$$n_{\bar{p}} = \alpha(M) n_\gamma f_{\bar{p}} \sqrt{\frac{m_{Pl}}{M}} = \frac{\alpha(M) n_\gamma f_{\bar{p}}}{(m_{Pl} t_e)^{\frac{1}{6}}} = \alpha f_{\bar{p}} \frac{10^{-7} n_\gamma}{t_e^{\frac{1}{6}}} . \tag{6.12}$$

Using (2.6) we have $n_\gamma = 2.5 \cdot 10^8 n_B / \Omega_B$, so that

$$n_{\bar{p}} = \frac{25 f_{\bar{p}} \alpha(M) n_B}{\Omega_B t_e^{\frac{1}{6}}} . \tag{6.13}$$

Since almost all the $\bar{p}$'s produced in PBH evaporation at the RD stage annihilate, one obtains from (6.13)

$$\bar{r} = \frac{n_{\bar{p}}}{n_B} = \frac{25\alpha(M)}{\Omega_B t_e^{\frac{1}{6}}} f_{\bar{p}} . \tag{6.14}$$

Equation (6.14) is true for PBHs evaporating at the RD stage. PBHs with mass $10^{13}\,\text{g} < M < 10^{15}\,\text{g}$ evaporate at the matter-dominated stage at $t > 10^{12}$ s. For such PBHs eq. (6.11) gives

$$n_{\bar{p}} = \frac{M_{\bar{p}}}{m_{Pl}^2} f_{\bar{p}} \alpha(M) \varrho_B = \frac{M_{\bar{p}} m_p}{m_{Pl}^2} f_{\bar{p}} \alpha(M) n_B , \tag{6.15}$$

so that the relative density of generated $\bar{p}$'s is equal to

$$\bar{r} = n_{\bar{p}}/n_B = \alpha(M) f_{\bar{p}} \cdot 10^{-1} . \tag{6.16}$$

Note that for a homogeneous distribution of matter $(n_B = \langle n_B \rangle)$ at the matter-dominated stage the magnitude (6.16) does not coincide with f—the fraction of annihilation antiprotons—since

$$t_e/\tau_{ann} = n_B (\sigma v)_{ann} t_e < 1 \tag{6.17}$$

at $t_e > t_{\bar{p}} = 2 \cdot 10^{14} \Omega_B$ s, *i.e.* the annihilation time scale is larger than the cosmological one, so that antiprotons do not manage to annihilate.

However, at the stage of galaxy formation $t > t_G \sim 10^{16}$ s neither matter, nor PBHs are distributed homogeneously.

PBHs are clustered in galaxies, and their evaporation provides a local (galactic) source of the antiproton component of cosmic rays.

Antiprotons evaporated in the period $t_{\bar{p}} < t < t_G$ experience a red-shift of their energy, but this red-shifting $\sim (t/t_G)^{\frac{4}{3}}$ may provide condensation in galaxies of only very slow $\bar{p}$'s (having initial energies $\sim 10^{-6} (t/t_G)^{-\frac{4}{3}}$ GeV). More energetic $\bar{p}$'s maintain isotropic background antiproton fluxes, distributed homogeneously in the Universe.

6·1.3. Antiprotons from PBHs and small-scale inhomogeneities. In the previous subsection the density of antiprotons produced in PBH evaporation was related to the contribution $\alpha(M)$ of PBHs with mass M to the total cosmological density at the moment t_e of their evaporation. The value of $\alpha(M)$ is determined by the spectrum of initial adiabatic fluctuations. In general, the relationship between $\alpha(M)$ and small-scale inhomogeneities is not simple, due to possible changes of the equation of state of the early Universe (see sect. **7** and **8**). However, for any given GUT model predicting these changes, such a relationship may be established. As a zeroth-order approximation we shall neglect the changes of equation of state in the early Universe. In this approximation we simply interpolate the equation of state $p = \varepsilon/3$, proven for $t > 1\,\text{s}$, within the first second of the cosmological expansion. This means that we assume that the dominance of ultrarelativistic particles and antiparticles, being in thermal equilibrium with the radiation, starts from the very beginning of the cosmological expansion, or, at least, from the Planck time. In this picture PBHs may be formed if, at the moment when the fluctuation encompasses the mass of the cosmological horizon, its amplitude is of order 1. For a given spectrum (6.1) the probability of PBH formation is given by the Gaussian distribution

$$W_{\text{PBH}}(M) \sim \exp[-1/\delta^2(M)]\,. \tag{6.18}$$

So, at the moment $t_f = (M/m_{\text{Pl}})t_{\text{Pl}}$ when the fluctuation of scale M enters the horizon (*), the fraction of matter $\beta(M)$ being within PBHs of mass M is given by

$$\beta(M) = \left.\frac{\varrho_{\text{PBH}}(M)}{\varrho_{\text{tot}}}\right|_{t=t_f} = W_{\text{PBH}}(M)\,. \tag{6.19}$$

At $t > t_f$ the relative contribution of PBHs to the total density increases, since $\varrho_{\text{PBH}}(M) \simeq M n_{\text{PBH}}(M)\,\varrho_{\text{tot}} \simeq 3Tn_r, n_{\text{PBH}}/n_r = \text{const}$ and

$$\frac{\varrho_{\text{PBH}}(M)}{\varrho_{\text{tot}}} \propto \frac{M}{T} \propto t^{\frac{1}{2}}\,. \tag{6.20}$$

PBHs with mass $M < M_{\text{RD}} \sim 10^{13}$ g evaporate before the end of the RD stage, so relationship (6.20) is valid all the time from the moment t_f of PBH formation up to the moment t_e of evaporation.

So for $M \leqslant M_{\text{RD}}$

$$\frac{\alpha(M)}{\beta(M)} = \left.\frac{\varrho_{\text{PBH}}(M)}{\varrho_{\text{tot}}}\right|_{t=t_e} \Big/ \left.\frac{\varrho_{\text{PBH}}(M)}{\varrho_{\text{tot}}}\right|_{t=t_f} = \left(\frac{t_e}{t_f}\right)^{\frac{1}{2}} = \sqrt{\frac{(M/m_{\text{Pl}})^3 t_{\text{Pl}}}{(M/m_{\text{Pl}})t_{\text{Pl}}}} = \frac{M}{m_{\text{Pl}}}\,. \tag{6.21}$$

(*) *I.e.* the wave-length of the fluctuation λ is equal to the cosmological horizon $l_h = ct$, so that the mass encompassed by the fluctuation—the scale—is equal to the mass being at the moment t within the horizon.

More massive PBHs, $M > M_{\mathrm{RD}}$, evaporate after the end of the RD stage, at $t > t_{\mathrm{RD}} \sim 10^{12}$ s at the matter-dominated stage, when $\varrho_{\mathrm{tot}} = m_{\mathrm{p}} n_{\mathrm{B}}$.

In this case the growth of the relative PBH density (6.20) is valid only until t_{RD}, and $\varrho_{\mathrm{PBH}}/\varrho_{\mathrm{tot}} = \mathrm{const}$ at $t > t_{\mathrm{RD}}$. So for $M \gtrsim M_{\mathrm{RD}}$

$$\frac{\alpha(M)}{\beta(M)} = \left(\frac{t_{\mathrm{RD}}}{t_{\mathrm{f}}}\right)^{\frac{1}{2}} = \left(\frac{m_{\mathrm{Pl}}}{M}\right)^{\frac{1}{2}} \left(\frac{t_{\mathrm{RD}}}{t_{\mathrm{Pl}}}\right)^{\frac{1}{2}}. \tag{6.22}$$

With due account for possible changes of the equation of states eqs. (6.21) and (6.22) are to be modified. Two main possibilities of such changes arise: *a*) early dustlike (*) $p = 0$ stages of super-heavy particles or mini BH dominance and *b*) inflationary stage of exponential expansion (see sect. **7**, **8**). In the first case all the time during the early $p = 0$ stage $\varrho_{\mathrm{PBH}}/\varrho = \mathrm{const}$, so that during this stage there is no enhancement of the relative contribution of PBHs to the cosmological density.

The effect of an early dustlike stage with beginning at t_0 and end at t_1 on the PBH concentration may be written as follows:

$$\frac{\alpha(M)}{\beta(M)} = \sqrt{\frac{t_0}{\min(t_0, t_{\mathrm{f}})}} \sqrt{\frac{\min(t_{\mathrm{e}}, t_{\mathrm{RD}})}{\max(t_{\mathrm{f}}, t_1)}}. \tag{6.23}$$

We see from (6.23) that the presence of early dustlike stages results in a decrease of the relative contribution of PBHs to the total density at the moment of their evaporation. Inflationary stages of exponential expansion reduce the magnitude $\alpha(M)/\beta(M)$ exponentially $\propto \exp[-3Ht_{\mathrm{inf}}]$ if $t_{\mathrm{f}} < t_{\mathrm{inf}}$, *i.e.* if formation of PBHs takes place before the end of an inflationary stage t_{inf}. However, additional mechanisms of black-hole formation are possible in cases *a*) and *b*) (see sect. **8**).

6'2. *Domains of antimatter.* – It has been already pointed out that entropy fluctuations of large amplitude cause regions with negative baryon charge $\langle B \rangle < 0$, *i.e.* antimatter domains. Since annihilation is possible on the boundaries of domains only, the matter-antimatter domain structure provides a much longer survival of antibaryons as compared to the case of their homogeneous distribution. The time scale of annihilation is determined by the scale of the domains. The greater the size of a domain, the longer is its annihilation time scale, and the longer is the period in which antibaryons are present in the Universe.

(*) Case *a*) implies the stage of dominance of nonrelativistic particles of mass m in the cosmological density. Their pressure p is of the order $p \sim m\gamma^2 n_{\mathrm{m}}$ being much smaller than the energy density $\varepsilon \sim mc^2 n_{\mathrm{m}}$: $p \ll \varepsilon$. So this stage of expansion may be considered as an expansion of dustlike matter with negligible ($p \approx 0$) pressure.

It is convenient to characterize the scale of a domain by the total baryon charge within it (*i.e.* by the total amount of antibaryons, or their excess within the domain). This quantity is conserved during the expansion, until the approximation of the thin boundary for annihilation is valid. The evolution of the matter-antimatter domain structure includes *a*) disappearance (annihilation) of small-scale domains, *b*) coalescence, *i.e.* formation of domains of larger scale from smaller-scale domains, and *c*) evolution of large-scale domains owing to annihilation in the thin layer of the matter-antimatter domain boundary.

Let us consider these effects, bearing in mind their possible relationship with the light-element abundances in the Universe (see sect. **4**). A detailed discussion of the evolution of the domain structures may be found in reviews [1, 6].

6.2.1. Annihilation of small-scale domains. Consider a domain of scale $N_{\bar{B}}$. Then for the baryon charge density $n_{\bar{B}}$ the size of the domain is of the order of

$$l \sim (N_{\bar{B}}/n_{\bar{B}})^{\frac{1}{3}} . \tag{6.24}$$

We may neglect in a first approximation the finite width of the boundary, treating it as infinitesimally thin. The annihilation of a domain is viewed in this approximation as a movement of this thin boundary with the speed u_b determined by the rate of annihilation. The latter is in its turn determined by the rate of baryon charge diffusion towards the boundary. Thus the time scale of annihilation of a domain is determined by the diffusion time scale.

At different stages of the cosmological expansion the diffusion of the baryon charge is determined by different processes.

At $t < 10^{-5}$ s gluons and $q\bar{q}$ pairs are in equilibrium with relativistic particles and radiation. The baryon charge of the antimatter domains is represented in this period by a small $(n_{\bar{B}}/n_\gamma < n_B/n_\gamma \sim 10^{-8} \div 10^{-10})$ $\bar{q}$ excess. The diffusion of this small excess towards the boundary determines the time scale of the domain annihilation. The coefficient of diffusion is given by

$$D = \tfrac{1}{3}\lambda v , \tag{6.25}$$

where v is the velocity and λ is the mean free path of the diffusing particle. In the considered period quarks are relativistic, so

$$v = c . \tag{6.26}$$

The mean free path λ for $\bar{q}$ migrating towards the boundary of a domain is determined by collisions with relativistic quarks, antiquarks and gluons, whose density is of the order of the density of radiation, *i.e.*

$$n_{\mathrm{q}} \sim n_{\bar{\mathrm{q}}} \sim n_{\mathrm{g}} \sim n_{\gamma} \sim T^3 . \tag{6.27}$$

The cross-section of these collisions may be estimated as (at $T \gtrsim 300$ MeV)

$$\sigma \sim \frac{\alpha_{\mathrm{c}}^2}{T^2} C , \tag{6.28}$$

where α_{c} is the QCD constant, and C is the colour screening factor in the colour plasma, similar to the Coulomb factor in the case of ordinary plasma. Combining (6.27) with (6.28), one obtains

$$\lambda = (n_{\mathrm{q}}\sigma_{\mathrm{q}\bar{\mathrm{q}}} + n_{\bar{\mathrm{q}}}\sigma_{\bar{\mathrm{q}}\bar{\mathrm{q}}} + n_{\mathrm{g}}\sigma_{\bar{\mathrm{q}}\mathrm{g}})^{-1} \sim (3n_{\gamma}\sigma)^{-1} \sim \frac{1}{3\alpha_{\mathrm{c}}^2 TC} . \tag{6.29}$$

If we take into account (6.25), (6.26) and (6.29), the coefficient of diffusion D is given by

$$D \simeq \frac{1}{9\alpha_{\mathrm{c}}^2 TC} \simeq \frac{10^{-4}}{T_{\mathrm{GeV}}} \frac{\mathrm{cm}^2}{\mathrm{s}} = \frac{3\cdot 10^9}{Z} \frac{\mathrm{cm}^2}{\mathrm{s}} . \tag{6.30}$$

Equation (6.30) is valid at red-shifts $Z > 10^{13}$. At the moment t $\bar{\mathrm{q}}$'s can migrate to the distance $l_Z = \sqrt{6Dt}$. At the RD stage

$$t = \frac{3\cdot 10^{19}}{Z^2}\ \mathrm{s} , \tag{6.31}$$

so $l_Z = 7\cdot 10^{14} Z^{-\frac{3}{2}}$ cm at the moment t corresponding to the red-shift Z. When l_Z is equal to the size of a domain, the domain dissipates, so the value of l_Z at $Z = 10^{13}$ ($t \sim 10^{-5}$ s) determines the maximal size of dissipated domains. To obtain the modern size of such domains, the value of l_Z is to be multiplied by the factor $1 + Z \approx Z$, accounting for the cosmological expansion

$$l_Z|_{\mathrm{modern}} = ZL_Z = 7\cdot 10^{14} Z^{-\frac{1}{2}}\ \mathrm{cm} . \tag{6.32}$$

Thus only domains of scale $N_{\bar{\mathrm{B}}}$ larger than

$$N_{\bar{\mathrm{B}}} \sim n_{\bar{\mathrm{B}}} l_Z^3 \sim \bar{r}\Omega_{\mathrm{B}} \cdot 10^{39} Z^{-\frac{3}{2}}|_{Z\simeq 10^{13}} = 3\cdot 10^{20} \bar{r}\Omega_{\mathrm{B}} \tag{6.33}$$

survive until $Z \sim 10^{13}$ ($t \sim 10^{-5}$ s). In (6.33) $\bar{r}$ is the ratio of antibaryon excess within the domain $\Delta n_{\bar{\mathrm{B}}}$ to the average excess n_{B} (in baryon-symmetric models $\bar{r} = 1$) and $\Omega_{\mathrm{B}} = \varrho_{\mathrm{B}}/\varrho_{\mathrm{c}}$ (see sect. **2**).

After $t \sim 10^{-5}$ s, when the temperature drops below ~ 300 MeV, coloured quarks and gluons combine into colourless hadrons (pions, nucleons, antinucleons), decays of pions and local annihilation of nucleons and antinucleons take place, so that by $t \sim 10^{-3}$ s (at $T \sim 20$ MeV) only the local baryon (antibaryon) excess is left. Until $t \sim 10^2$ s (*i.e.* up to the beginning of the stage of cosmological nucleosynthesis) this excess is represented by free neutrons and protons (and, accordingly, antinucleons in the antimatter domains). Electrons and positrons were in equilibrium with the radiation. At $T \gtrsim 1$ MeV they are relativistic, so that their equilibrium concentration is by $8 \div 10$ orders of magnitude higher than the nucleon (antinucleon) concentration. In this period the free antineutron's diffusion towards the boundary is the most essential mechanism of antimatter domain dissipation. At $T > 1$ MeV weak-interaction processes provide effective $\mathrm{n} \leftrightarrow \mathrm{p}$ conversion, so that the total baryon charge of a domain may have effectively migrated towards the boundary by the effect of antineutrons.

The diffusion rate of antineutrons is determined by their scattering on antinucleons and on electrons and positrons. The latter process, though having a very small cross-section, gives an essential contribution to the diffusion rate owing to the much larger concentration of e^+ and e^- as compared to the antinucleon one. So for the mean free path of antineutrons one obtains [6]

$$\lambda_{\bar{\mathrm{n}}} = (n_{\mathrm{e}} \sigma_{\bar{\mathrm{n}}\mathrm{e}} + n_{\bar{\mathrm{n}}} \sigma_{\bar{\mathrm{n}}\bar{\mathrm{n}}})^{-1} \tag{6.34}$$

and for the coefficient of diffusion

$$D_{\bar{\mathrm{n}}} \simeq \tfrac{1}{2} \tfrac{1}{3} \lambda_{\bar{\mathrm{n}}} v_{\bar{\mathrm{n}}} \,. \tag{6.35}$$

The additional factor $\frac{1}{2}$ is introduced into (6.35) since a nucleon spends, at $T > 1$ MeV, as a neutron half of its time. Cross-sections of $\bar{\mathrm{n}}\mathrm{e}$ and $\bar{\mathrm{n}}\mathcal{N}$ scatterings in (6.34) are given by [6]

$$\sigma_{\bar{\mathrm{n}}\mathrm{e}} \simeq 4 \cdot 10^{-37}\, T \,\mathrm{cm}^2 \qquad \text{and} \qquad \sigma_{\bar{\mathrm{n}}\mathcal{N}} \approx \frac{4 \cdot 10^{-24}}{T} \,\mathrm{cm}^2 \,, \tag{6.36}$$

where T is in MeV.

Recalling that $n_{\bar{\mathcal{N}}} = \bar{r} r_{\mathrm{B}} n_{\gamma}$ (where $\bar{r} = n_{\bar{\mathrm{B}}}/n_{\mathrm{B}}$ and $r_{\mathrm{B}} = n_{\mathrm{B}}/n_{\gamma}$) and $n_{\mathrm{e}} \approx \frac{3}{2} n_{\gamma}$, one obtains [6] from (6.35) and (6.36)

$$l_z \approx \frac{3 \cdot 10^4 \,\mathrm{cm}}{T^{\frac{11}{4}}} \left[1 + \frac{10^{8.8} \bar{r} r_{\mathrm{B}}}{T^2}\right]^{-\frac{1}{2}} , \tag{6.37}$$

where T is in MeV.

This gives for the maximal scale of dissipating domains (for $\bar{r} r_{\mathrm{B}} \ll 10^{-8.8}$)

$$N_{\bar{\mathrm{B}}} \sim \bar{r} r_{\mathrm{B}} n_{\gamma} l_z^3 \sim \bar{r} r_{\mathrm{B}} \frac{10^{45}}{T^{\frac{21}{4}}} \,. \tag{6.38}$$

So by the first second of the expansion (at $T \sim 1$ MeV) only domains with scale larger than $N_{\bar{\mathrm{B}}}^{(\bar{\mathrm{n}})} \sim \bar{r} r_{\mathrm{B}} \cdot 10^{45}$ have survived.

It was noted in [1] that at $t \leqslant 10^2$ s the presence of e^+e^- pairs in equilibrium with the radiation provides (owing to $n_{e^+} + n_{e^-} \gg n_{\bar{p}}$) a strong enhancement of the coefficient of proton diffusion (as compared to the case of equal concentrations of light and heavy charged particles $n_{e^+} = n_{\bar{p}}$ considered below). According to [1] diffusion of antiprotons in this period may result in the annihilation of domains on a scale smaller than

$$N_{\bar{B}} \sim 10^{62} \bar{r} r_B \,. \tag{6.39}$$

After local e^+e^- annihilation at $t \sim 10^2$ s the densities of light charged particles (electrons in the matter domains and positrons in the antimatter domains) are equal inside domains to the respective densities of heavy charged particles (p and $\bar{p}$, nuclei and antinuclei). In this case the diffusion rate is determined by the radiation friction of electrons and positrons—the radiation friction of heavy particles is negligible, but due to electrostatic forces they cannot migrate faster than light particles. The theory of diffusion at the RD stage is given in [1]. Its main result is that at $10^2\,\mathrm{s} \leqslant t \leqslant 10^{13}$ s

$$D = \frac{3ckT}{3\varepsilon_\gamma \sigma_T} = 0.6 \cdot 10^{32} Z^{-3} \frac{\mathrm{cm}^2}{\mathrm{s}}\,, \tag{6.40}$$

where

$$\sigma_T = \frac{8\pi}{3}\left(\frac{\alpha}{m_e}\right)^2 \simeq 6.65 \cdot 10^{-25}\ \mathrm{cm}^2$$

is the cross-section of Thomson photon-electron scattering, and ε_γ is the radiation energy density. So by the moment t, corresponding to the red-shift Z, the domains of a scale smaller than

$$N_{\bar{B}}(Z) = \frac{5 \cdot 10^{72}}{Z^{\frac{9}{2}}} \bar{r} r_B \tag{6.41}$$

must annihilate.

At $t = t_{rec} \sim 10^{13}$ s recombination takes place, so that neutral atoms are formed. It was suggested, however (cf. [38]), that the radiation from the regions of annihilation ionizes the nearly layers, so that expression (6.40) for the diffusion coefficient is valid for the periods $t > t_{rec}$ also. Taking formally relationship (6.41), one obtains for $Z \sim 1$ the mass of the survived domain:

$$M \sim 10^{15} \Omega_b \bar{r} M_\odot \,. \tag{6.42}$$

At $\bar{r} = 1$, *i.e.* in the baryon-symmetrical case, one obtains a mass of a domain rather close to the mass of a supercluster of galaxies. This formal coincidence with the scale of the cell structure of the Universe was used in [28] as an argument in favour of the baryon-symmetric model.

6'2.2. The effect of domain annihilation on nucleosynthesis. It was argued in [38] that small-scale domains coalesce into domains of large scale. The very physical picture of coalescence seems [1] not very reliable. It was noted in [1] that in the baryon-symmetric case [38] $\bar{r} = 1$ formation of domains of mass $\Omega_b \cdot 10^{11} M_\odot$ from small-scale domains results in a release of energy 20 times larger than the radiation energy density that is completely excluded by the restrictions on the possible distortions of the thermal-background spectrum. In the baryon-asymmetric case $\bar{r} \ll 1$, and coalescence cannot result in such a drastic contradiction with observations of electromagnetic radiation. However, if antimatter domains have characteristic scales $N_{\bar{B}}$, the main observational effect of their annihilation may be expected at the time when the scale of the domains is comparable with the diffusion scale l_Z. So in our further consideration of the limits on the matter-antimatter domain scales, determined by the parameters of GUTs (see sect. **8**), we shall restrict ourselves to the effect of annihilation in the period when $L = l_Z$. For the same reason we shall put aside the question on large-scale domain annihilation, since at $l_Z \ll L$ the value of f—the fraction of annihilated matter—decreases by the factor $(l_Z/L)^2$ (see [1]) as compared to the ratio $\bar{r}$ of the cosmological antibaryon and baryon densities. A negligible fraction of antiprotons annihilates, so the effect of annihilation on nucleosynthesis is negligible.

On the basis of the restrictions on f from observations of γ-radiation and of the thermal-radiation spectrum (see sect. **3**), the annihilation of domains of scale $N_{\bar{B}}$ (*)

$$(10^{45} \div 10^{62}) r_B \bar{r} < N_{\bar{B}} < 10^{67} r_B \bar{r} \tag{6.43}$$

seems to be of the utmost interest from the viewpoint of the effect of the annihilation on nucleosynthesis. These domains annihilate in the period $10^3\,\text{s} < t \leqslant t_{\text{rec}}$, for which the limits on the possible amount of annihilated antimatter, obtained from the measurements of the thermal-background spectrum, are rather weak (see subsect. **3**'3). Annihilation of domains (6.43) results in changes of light-element abundances. The possibility of formation of the domain structure on the scale (6.43) within the frame of GUTs is discussed in sect. **8**.

7. – Baryon-asymmetric model and the very early Universe.

The picture of the evolution of the Universe accepted at present (standard) is based on the statement that all the visible astronomical objects consist of matter, that there is no antimatter in the modern Universe. According to this

(*) The uncertainty in the lower limit is due to the difference of the estimates [1] and [6] of the scale of domains annihilated by the end of the nucleosynthesis.

picture, practically all the antiparticles, being in equilibrium with the particles and the radiation in the early Universe, must have annihilated in the course of the successive expansion (antibaryons after 10^{-3} s, positrons after 10^{2} s). The relatively small, as compared to the number of photons, number of baryons which are left after annihilation must have been given initially as an excess of baryons over antibaryons. The widespread prejudice against the baryon-asymmetrical model concerns that very baryon excess given initially. Its arbitrariness looked unesthetical, as compared to the esthetically attractive baryon-symmetrical cosmology. However, recently this theoretical « ugliness » of the baryon-asymmetrical cosmology was removed. It turned to be possible to relate this excess to the fundamental properties of particle physics.

7·1. *Baryon charge generation in the early Universe.* – The baryon excess is to be put into the Universe « by hand » from the very beginning, if the baryon number is conserved. However, the baryon charge (baryon number) differs substantially from the electric charge—there is no long-range field induced by the baryon charge, so its possible nonconservation would not result in a dramatic instantaneous change of this field.

No fundamental physical grounds prevent baryon nonconservation. Baryon charge conservation means only that in all the known reactions the number of baryons minus the number of antibaryons is conserved and that the lightest baryon—the proton—is stable. The latter is proven with high accuracy—the proton lifetime, as experiments show, must be greater than 10^{30} y. So, in fact, baryon conservation does not lie on any fundamental ground. It is simply an experimental fact, proved with high precision. SAKHAROV and KUZMIN were the first who pointed that baryon nonconservation may take place in particle interactions, inducing matter dominance in the modern Universe [45, 46]. They suggested that the processes between quarks and leptons (ℓ) of the type $\overline{\mathrm{dd}} \to \mathrm{d}\ell^+$, or $\mathrm{uu} \to \overline{\mathrm{d}}\ell^+$, or $\mathrm{ud} \to \overline{\mathrm{d}}\nu$ might have taken place in the early Universe, implying baryon charge generation. So, in the Universe, initially baryon symmetric, baryon asymmetry, *i.e.* net nonzero baryon charge, may arise owing to such processes. However, baryon charge generation implies additional conditions to be fulfilled. Baryon nonconservation only is not sufficient to produce baryon excess, since 1) the principle of detailed balance, being valid for any system with C and CP conservation, 2) thermodynamical equilibrium, implying that, in the absence of detailed balance, the rate of processes going from a given initial state to all the final states is equal to the sum of rates of the processes from all the possible states to a given initial state, preclude baryon asymmetry generation. Thus C and CP violation is to be evoked as well as inequilibrium conditions are to be realized. It turned out that all the three conditions may be fulfilled within the framework of GUT cosmology. A detailed discussion of the mechanisms of generation of baryon asymmetry in the early Universe may be found in reviews [5, 7, 8].

GUTs, being the extension of the gauge theory of unified weak and electromagnetic interactions and quantum chromodynamics, treat leptons and quarks on an equal footing, arranging them in the same representation of the underlying symmetry, *i.e.* they are considered as different states of one particle. All the interactions are considered then as local gauge transformations from one state to another state. The theory is based on the invariance of the interactions under such transformations, what leads inevitably to the existence of massless gauge vector bosons, mediating all the particle interactions. So all the forces turn to be long-range ones. But we know that weak and strong forces are of short range. To account for the observed difference in the range of the forces, the respective gauge bosons are to be made massive. The mechanism of generation of masses within the frame of local gauge theory was developed by HIGGS [47]. In this mechanism an auxiliary scalar field is introduced (see **7**'2.1) whose interactions with gauge bosons and fermions induce their masses. The simplest example of GUT is SU_5. GUT's interaction is induced by exchange of various gauge bosons. There are 24 intermediate bosons in SU_5. Twelve of them are those mediating electromagnetic, weak and strong interactions: 1 photon + 3 weak intermediate bosons + 8 gluons. There are twelve « new » interactions too, mediated by leptoquark X and Y bosons. X-bosons induce lepton-quark transitions $qq \to X \to \bar{q}\ell$. The mass of X-bosons is typically of order $M_X \sim (10^{15} \div 10^{15})$ GeV, so that they mediate proton decay $p \to e^+\pi^0$ or $p \to \pi^+\nu$ with a lifetime $\tau_p \gtrsim 10^{30}$ y. Leptoquarks cannot be produced neither in accelerators nor in cosmic rays, but they must have been present in the very early Universe, when the temperature was $T \gtrsim M_X$, *i.e.* at $t \lesssim 10^{-35}$ s. At these temperatures all the zoology of GUTs (superheavy Higgs mesons, superheavy leptons or quarks, ordinary Higgs and intermediate W-bosons, gluons, quarks, leptons, etc.) was present in equilibrium. All the kinds of processes, including the baryon-nonconserving ones, were possible. We shall restrict ourselves for simplicity only to X-bosons to give the main idea of baryon asymmetry generation.

Let us assume for simplicity that the leptoquark has only two modes of decay: $X \to qq$ (with branching ratio r) and $X \to \bar{q}\ell$ (with branching ratio $1 - r$). Then the corresponding antileptoquark $\bar{X}$ decays into $\bar{q}\bar{q}$ (branching ratio $\bar{r}$) and $q\bar{\ell}$ (branching ratio $1 - \bar{r}$). Owing to CPT, the lifetimes of X and its antiparticle $\bar{X}$ are to be equal. But if C and CP are violated, there is no detailed balance and the branching ratios of the respective modes are not equal, *i.e.* $r \neq \bar{r}$. So there is baryon number nonconservation and C (CP) violation in X decays. However, both these conditions are not sufficient for baryon excess production, since in the thermodynamical equilibrium there is detailed balance for direct and inverse reactions, so that the $X \to qq(\bar{q}\ell)$ decay and the inverse reaction $qq(\bar{q}\ell) \to X$ have equal rates. No net baryon excess arises in equilibrium. Nonequilibrium processes are needed. When the temperature falls below m_X, the concentration of X-bosons is « frozen », and their decay goes out of equilibrium. So, in X and $\overline{\mathrm{X}}$ decays a baryon excess

$$\Delta B = r\tfrac{2}{3} - \bar{r}\tfrac{2}{3} - \tfrac{1}{3}(1-r) + \tfrac{2}{3}(1-\bar{r}) = r - \bar{r}$$

per decay is generated.

Multiplying this excess by the X and $\overline{\mathrm{X}}$ density at the moment of decay n_{X} and dividing by the number density n_{r} of all the other relativistic particles being in equilibrium at the moment of decay, we obtain the magnitude of the baryon asymmetry

$$\frac{\Delta B}{n_\gamma} = (r - \bar{r})\frac{n_{\mathrm{X}}}{n_{\mathrm{r}}}. \tag{7.1}$$

The magnitude and the very sign of the asymmetry depend on the magnitude $r - \bar{r}$, which is in its turn determined by the sign of CP-violating phases. This point will be of great importance for future discussions (see sect. **8**).

7'2. *Physics of the very early Universe.* – Based on GUTs the picture of the first second of expansion may be analysed. Since the parameters of GUTs, as well as the correct GUT itself, are not established now, this picture is ambiguous.

Besides the above-mentioned possibility of baryon charge generation, GUTs predict a number of nontrivial cosmological consequences: phase transitions, production of magnetic monopoles and of some other new particles, etc. Some GUT models predict matter-antimatter domain structure, massive walls and strings, heavy particles or anomalous vacuum dominance stages. Here we shall make a brief comment, more detailed considerations concerning possible sources of antinucleons induced by GUTs will be given in sect. **8**.

7'2.1. Phase transition in the early Universe. The modern approach to the description of the difference of the fundamental forces implies the mechanisms of spontaneous breakdown of the underlying gauge symmetry in close analogy to the theory of superfluidity, superconductivity or ferromagnetism. The interactions possess the symmetry of the theory, however the ground state (vacuum) of the theory is asymmetrical, inducing symmetry breaking. The difference between the observed properties of the fundamental interactions is ascribed to the existence of an auxiliary scalar Higgs field. Its self-interactions make the state with nonzero vacuum expectation value of this field energetically favourable. Interactions of fermions and gauge bosons with the condensate induce their masses. So, owing to Higgs condensate, symmetry breaking occurs.

Kirzhnitz and Linde [18] have pointed out that in strict analogy with ferromagnetism (or other similar physical phenomena related to spontaneous symmetry breaking) at high temperatures restoration of the symmetry must take place: a certain critical temperature exists over which condensate and, consequently, fermion and boson masses disappear, as ferromagnetic properties disappear above the Curie-Weiss temperature. At such high temperatures the ground state with nonzero vacuum expectation value of the Higgs field turns out to be energetically unfavourable, as compared to the state with zero field. Symmetry is restored since the ground state is symmetrical at high temperatures. So, in the beginning of the expansion, the symmetrical phase is realized, when all the particles are massless, when all the interactions are unified.

In the course of the expansion the temperature decreases, so when it falls below the critical temperature, a phase transition to the asymmetrical state must take place. Within the frame of GUTs there must have been at least two phase transitions in the early Universe: the transition to the phase in which the unified asymmetry of GUT is broken, so that strong and unified electroweak interactions are separated, and the Weinberg-Salam phase transition from the unified electroweak interaction symmetry to the phase in which only electromagnetic gauge symmetry remains unbroken, in which W and Z bosons acquire mass owing to the interaction with the Higgs condensate, so that weak and electromagnetic interactions are separated. The history of phase transitions in the early Universe may be even more complicated, since the structure of Higgs interactions may induce phase transitions to some intermediate phases. Phase transitions of another type may also occur in the early Universe—the transition from the unconfined to the confined phase of quarks and gluons, induced by the confining forces of quantum chromodynamics.

To make clear the idea of symmetry restoration, consider the model Lagrangian (in fact, the Lagrangian density) for the scalar field φ with parameters m and λ

$$\mathscr{L} = \frac{1}{2}\left(\frac{\partial\varphi}{\partial x_\mu}\right)^2 - \frac{1}{2}\, m^2\varphi^2 - \frac{1}{4}\,\lambda^2\varphi^4 \tag{7.2}$$

and the Hamiltonian (Hamiltonian density)

$$H = \frac{1}{2}\left(\frac{\partial\varphi}{\partial t}\right)^2 + \frac{1}{2}\left(\frac{\partial\varphi}{\partial \boldsymbol{x}}\right)^2 + \frac{1}{2}\, m^2\varphi^2 + \frac{1}{4}\,\lambda^2\varphi^4\,. \tag{7.3}$$

For a constant field $\varphi(\boldsymbol{x}, t)$

$$H = V(\varphi) = \tfrac{1}{2} m^2\varphi^2 + \tfrac{1}{4}\lambda^2\varphi^4\,. \tag{7.4}$$

We see that $H_{\min}$ corresponds to $\varphi = 0$, *i.e.* in the ground state (in the vacuum of the theory) no field is present. Consider now Lagrangian (7.2) with the « wrong » sign of the $m^2\varphi^2$ term:

$$\mathscr{L} = \frac{1}{2}\left(\frac{\partial\varphi}{\partial x_\mu}\right)^2 + \frac{1}{2}\, m^2\varphi^2 - \frac{1}{4}\,\lambda^2\varphi^4\,. \tag{7.5}$$

In this case for $\varphi = \text{const}$ the Hamiltonian

$$H = V(\varphi) = -\tfrac{1}{2}m^2\varphi^2 + \tfrac{1}{4}\lambda^2\varphi^4 \tag{7.6}$$

has a minimum at $\varphi = \pm m/\lambda$, and at $\varphi = 0$ H has a local maximum

$$\Delta H(\varphi = 0) = H(\varphi = 0) - H_{\min} = \frac{1}{4}\frac{m^4}{\lambda^2}.$$

The state without mean field is energetically unfavourable. In the ground state (vacuum) a nonzero field φ is present. If there is local gauge symmetry (in our simple example $\varphi \to \varphi \exp[i\chi(x)]$, $\chi(x)$ is the local gauge phase), there is a gauge field $\mathscr{A}_\mu(x) \to A_\mu(x) + \partial_\mu \chi(x)$ and the gauge-invariant Lagrangian is to be written as

$$\mathscr{L} = [(\partial_\mu - A_\mu)\varphi]^2 - \tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} + \tfrac{1}{2}m^2|\varphi|^2 - \tfrac{1}{4}\lambda^2|\varphi|^4, \tag{7.7}$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$.

But, owing to the presence of the mean field $|\varphi| = m/\lambda$ in the vacuum of the theory, we have ($\varphi = m/\lambda = \text{const}$)

$$\mathscr{L}' = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{m^2}{\lambda^2}A_\mu A^\mu. \tag{7.8}$$

The Lagrangian (7.8) is no longer gauge invariant: its second term is not invariant with respect to the transformations $A_\mu \to A_\mu + \partial_\mu \chi$. Owing to the interaction with the mean field $\langle\varphi\rangle = m/\lambda$ vector bosons acquire nonzero mass.

Starting from the gauge-invariant Lagrangian, we come to the gauge-noninvariant Lagrangian (7.8). The interaction with the field φ induces the spontaneous breakdown of the underlying gauge symmetry. The Hamiltonian (7.6) is valid for zero temperature. At high temperatures ($T \gg m$) thermal fluctuations arise, inducing in the Hamiltonian the additional term $\sim \tfrac{1}{2}cT^2\varphi^2$ (c is constant). It leads to the Hamiltonian

$$H(T) = V(\varphi, T) = \tfrac{1}{2}(cT^2 - m^2)\varphi^2 + \tfrac{1}{4}\lambda^2\varphi^4. \tag{7.9}$$

We see that, owing to thermal fluctuations, the minimum of $H(T, \varphi)$ at $T > m/\sqrt{c}$ corresponds to $\varphi = 0$. So the symmetry is restored, no mean field is present in the ground state. Vector bosons have no mass. With the decrease of the temperature to $T \sim m/\sqrt{c}$ a phase transition from the disordered (symmetrical) phase to the ordered (asymmetrical) phase takes place.

7.2.2. Domain walls. Note that in our example $V(\varphi)$ possesses an additional discrete symmetry $\varphi \to -\varphi$. In the expanding Universe before phase transitions, the symmetry is restored, and the sign of φ is not fixed. This sign may change in space with time, owing to thermal fluctuations. The changes are correlated to the distances determined by the characteristic value (correlation length) of this fluctuation $l \sim 1/T$. So, when the phase transition takes place, the distribution of signs of φ is frozen, the uncorrelated phase may obtain different signs in different regions—domain structure arises. Since in in each domain a definite sign of φ is taken, the underlying discrete symmetry is broken in each domain. On the boundaries between domains with opposite signs the value of φ changes, as can be easily estimated from eq. (7.3), from $-m/\lambda$ to $+m/\lambda$. Owing to this change, a massive wall with mass on unit area $\sim \lambda(m/\lambda)^3$ and width $\sim (1/\lambda)(m/\lambda)$ arises. The presence of these walls in the Universe results in dramatic consequences—their existence in the modern Universe would have induced strong inhomogeneities at large scales, their presence at earlier stages must have modified the picture of the cosmological expansion to a dramatic extent. Soon after the phase transition, walls must have dominated in the cosmological density, so that at present their density must have been by many orders of magnitude higher than the matter density in drastic contradiction with the observations.

A physical example of discrete symmetry of this kind is *CP*. Thus the domain wall problem must arise in the models of « soft » *CP* violation, in which *CP* violation is induced after the phase transition owing to the nonzero *CP*-violating phase of the mean field. The uncorrelated sign of this phase before the phase transition must have led to domains of opposite sign of *CP* violation after the transition with massive walls on the boundaries. However (see sect. **8**), it may be possible to avoid the troubles of massive domain walls in the framework of « refined » GUTs.

7'2.3. GUT picture of the very early Universe. Within the frame of GUTs, it is possible in principle to answer the question on the equation of state of the very early Universe. In first approximation the answer is that at very high temperatures all the particles are relativistic, they are in equilibrium with the radiation, and thus the equation of state of the ultrarelativistic gas $p = \varepsilon/3$ (where p is the pressure and ε is the energy density) may be taken. This means that the law of expansion $T \propto t^{-\frac{1}{2}}$ proved for the stages after the first second may be extrpolated to the first second of the expansion. However, a detailed analysis of various GUT models shows that the picture is much more complicated. In fig. 2.1 the time scale of the very early Universe is given.

Several characteristic moments are predicted by the simplest version (SU_5) of GUTs. They are the following: 1) The phase transition from the SU_5 symmetric phase to the $SU_2 \times SU_2 \times U_1$ phase in which strong interactions are decoupled from electroweak interactions. $t_{\text{GUT}} \sim 10^{-35}$ s and the corresponding critical temperature is $T_c \sim M_X \sim 10^{15}$ GeV. The « soft » *CP* violation appearing in refined GUTs with a more complicated structure of phase transitions (see sect. **8**) takes place in this period. Soon after this moment baryosynthesis takes place. 2) The phase transition from the $SU_2 \times U_1$ symmetrical phase of unified electroweak interactions to the $U_{1,\text{e.m.}}$ phase of short-range (W and Z bosons acquire mass) interactions, $t_{\text{WS}} \sim 10^{-11}$ s. The critical temperature is $T_c \sim \Lambda_{\text{WS}} \sim 300$ GeV. 3) Confinement of colour. The transition from the quark-gluon plasma to the gas of hadrons, $t_c \sim 10^{-5}$ s. The critical temperature is $T \sim \lambda_{\text{QCD}} \sim 300$ MeV.

In SU_5 the « gauge desert » is predicted from Λ_{WS} up to M_X—no principally new physics arises in this energy interval. However, even within the frame of SU_5 the « gauge desert » does not necessarily imply constancy in the conditions of expansion from the GUT phase transition up to the WS phase transition. The problem of relic magnetic monopoles is to be mentioned here.

7·2.4. The problem of magnetic monopoles. An inevitable prediction of SU_5, as well as of all the other GUTs [48, 49], is the existence of isolated magnetic poles—magnetic monopoles [50, 51]. Their existence is the price to pay for the unification of the fundamental forces. The mass of monopoles is predicted [52] of the order of $m \geqslant 10^{16}$ GeV. When the temperature dropped below $T_0 = (n_m/3n_r)m \sim 10^7$ GeV, monopoles with frozen concentration [53, 54] n_m (relative to n_r) must have dominated in the cosmological density, so that the dustlike stage of monopole dominance with the equation of state $p = 0$ must have started. At this stage monopole density perturbations grow and the structure of monopole and antimonopole inhomogeneities evolves. Separated in the inhomogeneities from the cosmological expansion, monopoles and antimonopoles annihilate, so that owing to annihilation the dustlike stage of monopole dominance ends, and ultrarelativistic products of annihilation maintain the equation of state $p = \varepsilon/3$. However, the gravitational clumping of monopoles is not complete (more than 1 percent of them does not participate in the clumping). This fraction of monopoles left from the first stage arranges in the course of the successive expansion a new stage of monopole dominance, etc. It was shown in [55-57] that the scenario of successive monopole dominance stages cannot reduce the monopole concentration down to the existing observational upper limits. A strong contradiction with observations arises, being a serious problem for cosmological applications of GUTs. There are different approaches to avoid the cosmological overproduction of monopoles [55-60] in any way modifying the simple scheme of fig. 2.1.

7·2.5. Inflation of scales in GUTs. One possible solution of the monopole problem is related to the dynamics of GUT phase transitions. As was pointed out earlier (see subsect. **7·2.1**), the phase transition begins when the temperature drops below the critical one: $T \leqslant T_c \leqslant M_X$. But the transition to the new vacuum proceeds as a phase transition of the first order: the transition proceeds through the formation of bubbles of the « new » vacuum. If the corresponding nucleation rate, *i.e.* the rate of bubble formation, is low, the transition is suppressed. In this case the so-called « supercooling » takes place. The cosmological expansion goes beneath the critical temperature with energetically unfavourable « symmetrical » vacuum. This vacuum has very high energy density $\varrho_v \sim T_c^4$, dominating over the energy density of relativistic plasma and radiation $\varrho_r \sim T^4$ at $T < T_c$.

The relativistic invariance of this vacuum implies the equation of state $p = -\varepsilon$, *i.e.* the anomalous (symmetrical) vacuum is the medium with negative pressure. Owing to this equation of state, the essential property of the stage of anomalous-vacuum dominance (AVD) is the exponential inflation of all the scales. However, as all the states with negative pressure, the AVD state is unstable and the time scale of its instability determines the degree of the inflation of the scales. This time scale is determined by the nucleation rate of formation of bubbles of the « true » vacuum, as well as by the kinetics of the transition. With the decrease of the temperature the nucleation rate grows, and at a certain temperature, T_b, intensive bubble formation begins. Successive thermalization leads to the reheating of the Universe up to a temperature of the order $\sim T_c$—after thermalization all the energy density of the anomalous vacuum transforms into the energy density of ultrarelativistic plasma and radiation.

Owing to the inflation of scales at the AVD stage monopole production in the phase transition may be highly suppressed [61, 62]. Indeed, monopoles should be generated at these points, by the condition that the phase of normal vacuum is changed of $2\pi n$ along the closed circle around this point. If the above condition is fulfilled, the continuity of the field φ implies $\varphi = 0$ at these points, so it is the topology that provides the existence of points of anomalous vacuum—magnetic monopoles. The concentration of produced monopoles is determined by the density of such points, being determined in its turn by the mean distance over which the phase of the mean field is uncorrelated and thus may change essentially. This distance is of the order of the correlation lenght l_c for fluctuations of the mean field φ, so that the number density of produced monopoles is of the order $n_M \sim (1/l_c)^3$. The magnitude l_c is $l_c \sim 1/T_c$, so that

$$n_M \sim T_c^3 . \tag{7.10}$$

If the transition proceeds at $T_b \ll T_c$, the correlation length l_b is much larger than l_c, $l_b = l_c(T_c/T_b)$, so that, after the phase transition, the density of produced monopoles is much lower than (7.10):

$$n_M \sim (1/l_b)^3 \sim \left(\frac{T_b}{T_c}\right)^3 \left(\frac{1}{l_c}\right)^3 \sim \left(\frac{T_b}{T_c}\right)^3 T_c^3 \sim T_b^3 \ll T_c^3 . \tag{7.11}$$

One obtains the suppression of monopole production.

There are some other fundamental problems for which the inflational scenario may provide a solution [14, 61, 62]. The first problem is the horizon problem [14, 61, 62]: the presently observed part of the Universe is quite homogeneous and isotropic at very large scales, that is exhibits a remarkable similarity of conditions in different regions. These regions, being now within the modern cosmological horizon, were causally disconnected at an earlier stage of the expansion. So, their initial conditions should have been uncorrelated, making their similarity mysterious. The second problem is the « flatness problem »—why is the Universe so close to the flat one (*i.e.* the observed cosmological density, within an order of magnitude, is close to the critical density)? As was shown in [61], this condition is fulfilled for $\Omega \neq 1$ only in the case of extremely fine « tuning » between the main cosmological parameters at $t \sim t_{\mathrm{Pl}}$. These problems were hoped to be solved. However, in the model considered in [61], such a small T_{b} (~ 1 K) is necessary that successive thermalization is impossible up to the present time, in strong contradiction with observations. A specific model, avoiding the trouble of very small T_{b}, is actively developed by LINDE [62]. In this model, the transition to the asymmetric phase is not connected with the decrease of the temperature. It is induced by strong coupling of fermions (the effect of GUT interactions similar to confinement of quarks in quantum chromodynamics) and, owing to the specific form of the Higgs interaction potential, a delay of the exponential inflation is possible after the transition to the asymmetric vacuum. In this model the inflation of scales is given not by

$$l_{\mathrm{b}} \sim \frac{T_{\mathrm{c}}}{T_{\mathrm{b}}} l_{\mathrm{c}}, \tag{7.12}$$

but by

$$l_{\mathrm{b}} \sim l_{\mathrm{c}} \frac{T_{\mathrm{c}}}{\Lambda} \exp\left[Ht_{\mathrm{inf}}\right] \sim \frac{T_{\mathrm{c}}}{\Lambda} l_{\mathrm{c}} \exp\left[\frac{3H^2}{\Lambda^2}\right] \tag{7.13}$$

with $H \sim T_{\mathrm{c}}^2/m_{\mathrm{Pl}}$. Here Λ is the scale of the strong-coupling limit in GUTs.

Whatever possibility is realized, we may conclude that inflation of scales is possible within the frame of GUTs. It will be of great importance for some sources of antimatter in the Universe.

8. – Sources of antimatter in GUTs.

On the basis of GUTs new possibilities of formation of domains of antimatter and PBHs arise.

8.1. *Domains of antimatter and phase transitions in refined* GUTs. – The relationship between the baryon asymmetry of the Universe and baryon-nonconserving *CP*-violating inequilibrium processes in the early Universe provides after slight modifications the possibility of antimatter domain formation. Indeed (see sect. 7), the sign of the excessive baryon charge is determined by the sign of the *CP*-violating phase. If for some reasons in the period of baryosynthesis regions with opposite signs of this phase exist, domains of antibaryon excess arise.

Within the frame of GUTs two mechanisms of *CP* violation are possible: *a*) Hard *CP* violation, induced by *CP*-violating terms in Yukawa interactions of fermions with Higgs bosons. These terms are contained in the Lagrangian of the theory owing to the complexity of the corresponding constants and thus are given initially possessing no spatial variation of *CP*-violating effects in the Universe. *b*) « Soft » *CP* violation. In this mechanism first suggested by LEE [63] there are no *CP*-violating terms before the spontaneous breaking of the GUT symmetry, but the mean field of the condensate is complex, inducing a nonzero *CP*-violating effect. *CP* invariance is a discrete symmetry, so soft *CP* violation is an example of *spontaneous* breakdown of discrete symmetry.

In the case of soft *CP* violation the energies of vacua of different signs of *CP* violation phases are equal, *i.e.* there is degeneracy of vacuum states. Both signs of *CP*-violating phases are possible, both are realized after the phase transition in spatially separated regions. Opposite-sign domains with massive walls on the boundaries are to be formed.

The problem of the cosmological consequences of « soft » *CP* violation pointed out in [64] is related to the evident contradiction between the existence of massive walls within the cosmological horizon and observations.

However, it was shown recently in [65] that the problem of domain walls may be resolved. The baryon excess is created in a rather short period of cosmological expansion. It is sufficient for the formation of antimatter domains that « soft » *CP* violation be switched on during a limited time interval, including that short period. If afterwards the « soft » mechanism of *CP* violation is switched out, and the corresponding symmetry is restored, no domain walls originated by soft *CP* violation survive. No contradiction with the observational data arises. KUZMIN, SHAPOSHNIKOV and TKACHEV have shown that this scenario may be realized within the frame of SU_5 with enlarged Higgs sector. If there are several Higgs fields, it is possible to arrange their interactions in such a way that spontaneous *CP* violation takes place in a certain interval of temperatures, vanishing at high and low temperatures.

To illustrate their idea, let us return to the Hamiltonian (7.4) of the $\varphi_1 =$ const scalar field in which interactions with other constant Higgs fields $\varphi_2 \neq 0$, $\varphi_3 \neq 0$ are taken into account:

$$H = V(\varphi_1) = \tfrac{1}{2} m^2 \varphi_1^2 + \tfrac{1}{4}\lambda^2 \varphi_1^4 + \tfrac{1}{2}\lambda_{12}\varphi_1^2\varphi_2^2 + \tfrac{1}{2}\lambda_{13}\varphi_1^2\varphi_3^2 = \tfrac{1}{2}\tilde{m}^2\varphi_1^2 + \tfrac{1}{4}\lambda^2\varphi_1^4 . \tag{8.1}$$

The minimum of this potential corresponds to $\varphi_1 = 0$, if

$$\tilde{m}^2 + m^2 + \lambda_{12}\varphi_2^2 + \lambda_{13}\varphi_3^2 > 0 . \tag{8.2}$$

So we have the theory without condensate $\varphi_1 \neq 0$ at low temperatures. Thus the symmetry of the theory corresponding to φ_1 is not violated.

At very high temperatures thermal fluctuations are to be taken into account:

$$H(T, \varphi_1) = \tfrac{1}{2}(c_1 T^2 + c_2 T^2 + c_3 T^2 + \tilde{m}^2)\varphi_1^2 + \tfrac{1}{4}\lambda^2\varphi_1^4\,, \tag{8.3}$$

where the c_1-term is induced by thermal fluctuations of φ_1 owing to $\lambda^2\varphi_1^4$ interactions and thus $c_1 > 0$, c_2 and c_3 are induced by thermal fluctuations of φ_2 and φ_3 because $\lambda_{12}\varphi_1^2\varphi_2^2$ and $\lambda_{13}\varphi_1^2\varphi_3^2$ may be negative. We take $c_2 < 0$, $c_2 + c_1 < 0$ and $c_1 + c_2 + c_3 > 0$. There is no condensate of φ_1 at $T \gg \tilde{m}/\sqrt{c_1 + c_2 + c_3}$ (the minimum of $H(T, \varphi_1)$ is reached at $\varphi_1 = 0$), *i.e.* there is no violation of symmetry at very high temperature. But with the decrease of the temperature, at $T < m_3$, thermal fluctuations of φ_3 are suppressed, so that instead of (8.3) we have

$$H(T < m_3, \varphi_1) = \tfrac{1}{2}(c_1 T^2 + c_2 T^2 + \tilde{m}^2)\varphi_1^2 + \tfrac{1}{4}\lambda^2\varphi_1^4\,. \tag{8.4}$$

Since $c_1 + c_2 < 0$ and provided that $|(c_1 + c_2)m_3^2| > \tilde{m}$, an interval of temperatures $\tilde{m}/|c_1 + c_2|^{\frac{1}{2}} < T < m_3$ exists in which the minimum of $H(T < m_3, \varphi_1)$ is realized at

$$\varphi_1 = \frac{\sqrt{|c_1 T^2 + c_2 T^2 + \tilde{m}^2|}}{\lambda}\,,$$

i.e. a nonzero mean field φ_1 arises, inducing violation of the symmetry.

It was shown in [65] that a rather natural set of parameters of scalar-field interactions may provide the conditions pointed out in the given example. So it is possible to avoid the problem of relic domain walls. But for antimatter domains to survive until nucleosynthesis, their size is to be sufficiently large:

$$N_{\bar{\mathrm{B}}} > 10^{54}\bar{r}\Omega_{\mathrm{B}} = 3\cdot 10^{62} r_{\mathrm{B}}\bar{r}\,,$$

where $r_{\mathrm{B}} = n_{\mathrm{B}}/n_{\gamma}$, $\bar{r} = \Delta n_{\bar{\mathrm{B}}}/\Delta n_{\mathrm{B}}$.

At $t < 0.1$ s the size of these domains $l \sim (N_{\bar{\mathrm{B}}}/\bar{r} r_{\mathrm{B}} n_{\gamma})^{\frac{1}{3}}$ exceeds the cosmological horizon $l_{\mathrm{h}} \sim ct$. Another problem arises: the problem of domain scales.

This problem may be solved on the basis of a possible inflation of the scales in the phase transitions. However, the situation is not so simple. The AVD stage in the course of a phase transition may provide inflation of the correlation length (see (7.12))

$$l = l_0 \frac{T_{\mathrm{c}}}{T_{\mathrm{b}}}\,,$$

so that the final scale of the domains is of the order

$$N_{\bar{\mathrm{B}}} = \bar{r} r_{\mathrm{B}} T_{\mathrm{c}}^3 l_0^3 \left(\frac{T_{\mathrm{c}}}{T_{\mathrm{b}}}\right)^3. \tag{8.5}$$

Taking $T_c \sim 10^{15}$ GeV, $T_b \sim 10^8$ GeV, and $l_0 \sim 1/T_c$, one obtains

$$N_{\bar{B}} = \bar{r}r_B \cdot 10^{21}, \tag{8.6}$$

that is obviously insufficient to achieve the desirable scale $N_{\bar{B}} \sim \bar{r}r_B 3\cdot 10^{62}$. However, KUZMIN, SHAPOSHNIKOV and TKACHEV [66] have shown that even in the simplest case of SU_5 the picture of spontaneous breakdown of GUT symmetry down to the $SU_3 \times SU_2 \times SU_1$ symmetry of separated strong and electroweak interactions may be rather complicated, so that for a reasonable choice of the parameters of Higgs interactions intermediate symmetries and hence intermediate phase transitions take place. If in each such transition the AVD stage is realized, the scale of domains increases by the factor $(T_c/T_b)^3$ in each transition:

$$N_{\bar{B}} = \bar{r}r_B T_{c_1}^3 l_0^3 \left(\frac{T_{c_1}}{T_{b_1}}\right)^3 \left(\frac{T_{c_2}}{T_{b_2}}\right)^3 \left(\frac{T_{c_3}}{T_{b_3}}\right)^3. \tag{8.7}$$

So, to achieve the desirable scale (for $T_{c_i} \sim 10^{15}$ GeV and $T_{b_i} \sim 10^8$ GeV), at least two additional phase transitions must occur. As was shown in [67], there does exist the possibility of a chain of transitions:

$$SU_5 \to SU_4 \to SU_3 \times U_1 \to SU_3 \times SU_2 \times U_1,$$

or

$$SU_5 \to SU_4 \times U_1 \to SU_4 \to SU_3 \times SU_2 \times U_1.$$

So two additional transitions may arise rather naturally. To achieve larger scales (*i.e.* cluster of galaxies), more transitions are needed. However, within the frame of the model [62] suggested by LINDE, one transition is sufficient to obtain any desirable scale, owing to the exponential inflation of scales (see (7.13)).

In this case the scale of domains obtained is of the order

$$N_{\bar{B}} \sim \bar{r}r_B T_c^3 l_0^3 \left(\frac{T_c}{\Lambda}\right)^3 \exp\left[\frac{9T_c^4}{m_{Pl}^2 \Lambda^2}\right], \tag{8.8}$$

where $(T_c/\Lambda)^3$ is induced by the AVD stage, preceding the beginning of the transition at $T \sim \Lambda$.

Let us summarize the main points of antimatter domain formation in GUTs. There are two mechanisms of CP violation—a hard one and a soft one—and two sets of CP-violating phases φ_h and φ_s, respectively. When soft CP violation is switched on, domains of $\varphi_h + \varphi_s$ and $\varphi_h - \varphi_s$ arise with massive walls on the boundaries. Owing to succession of phase transitions or exponential inflation the scales of domains grow as (7.12) or (7.13). In both cases, when the final (the only) transition to the $SU_3 \times SU_2 \times U_1$ phase ends, in the subsequent reheating baryosynthesis starts inducing the baryon excesses $r_h + r_s$ and $r_h - r_s$

in the domains $\varphi_h + \varphi_s$ and $\varphi_h - \varphi_s$. If $r_h < r_s$, the latter domains are the domains with antibaryon excess. In the course of expansion after baryosynthesis soft *CP* violation is switched out, symmetry restoration occurs, so that at lower temperatures $\varphi_s = 0$ and massive walls disappear, but the baryon excess domain structure is retained, resulting after local annihilation at $T \leqslant 1$ GeV in matter-antimatter domain formation. The scale of antimatter domains is given by (8.7) or (8.8). For annihilated domains of given scale the relative amount of annihilated antibaryons is

$$\bar{r} = \frac{r_s - r_h}{r_s + r_h}.$$

Note that the suggested scenario revives the theory of isothermal (entropy) fluctuations (see sect. **6**). For an arbitrary relationship between r_h and r_s, baryon charge inhomogeneities are predicted and it is possible, in principle, to obtain any desirable scale of such inhomogeneities. However, in this scenario the spectrum of entropy fluctuations is determined by GUTs, so that a close relationship between particle physics and galaxy formation theory arises.

8'2. *Heavy metastable particles in* GUTs. – The possibility of matter-antimatter domains considered in the previous subsection is connected within the frame of GUTs with Higgs mechanism of GUT symmetry breakdown. As a result of such a breakdown massive particles arise with masses of the order of $(10^{12} \div 10^{16})$ GeV. Most of them have lifetimes $\tau(m)$ comparable with the respective cosmological time scale m_{Pl}/m^2, so that they cannot survive long after the temperature T drops down to $T \leqslant m$. This is the case for X-bosons, Higgs mesons and most heavy fermions, predicted by GUTs. However, there are special cases in which particles have a lifetime (if any) much greater than the cosmological time scale m_{Pl}/m^2 and thus may be called metastable. These particles can survive long after $T \sim m$. At $T \ll m$ they go out of equilibrium and their annihilation is frozen exceeding substantially the equilibrium ($\sim \exp[-T/m]$) one.

If the lifetime of the particles is greater than ~ 1 s, their presence in the Universe might have affected physical processes in the Universe after nucleosynthesis and, in particular, primordial abundances of light elements. Particles with a lifetime smaller than 1 s cannot influence directly the big-bang nucleosynthesis, since they decay earlier than thermonuclear reactions start. However, they might have dominated before decay in the cosmological density, thus providing the early « dustlike » stage of their dominance, at which PBH formation (see sect. 6) from small density perturbations is possible.

PBHs survive long after the decay of the particles they are originated from, thus providing indirect influence of such particles on the physics of expansion after 1 s. Directly or indirectly heavy metastable particles induce late antibaryon production owing either to decays of metastable particles (with $\tau > 1$ s) in $p\bar{p}$+anything, or to evaporation of PBHs (with $M < 10^{13}$ g). Let us discuss the relationship between the heavy metastable particles and PBHs.

8'2.1. Frozen concentration of metastable particles. The net effect of heavy metastable particles is determined by their frozen concentration. The picture of freezing of these particles is, in general, similar to the freezing of antinucleons (see sect. **2, 5**). A detailed consideration of the kinetics of freezing, based on the kinetic equation, may be found in ref. [1, 2, 5]. Our aim is to illustrate the principal idea by some examples.

1) Very heavy metastable quarks G are predicted by some asymptotically free models of SU_5 [68]. These quarks have exotic colour properties (they belong to te octet of $SU_{3,c}$), implying their stability relative to decays induced by weak, strong or electromagnetic interactions. They may be absolutely stable. However, processes induced by exchanges of superheavy X-bosons may induce their decay $G \to 2q\ell$ or 3q, where q are « ordinary » quarks. Their lifetime relative to this decay may be estimated as

$$\tau_G \sim \frac{M_X^2 192\pi^3}{g_c^4 N_{ch} m_G^5 a^2}, \tag{8.9}$$

where M_X, m_G are masses of X-bosons and heavy quarks (respectively), g is the gauge constant of SU_5, N_{ch} is the number of decay modes and a is the mixing angle characterizing the $G \to q$ transition; τ_G exceeds $t \sim m_{Pl}/m_G^2$ for

$$m_G < \left(\frac{M_X^4 192\pi^3}{g^4 m_{Pl} a^2 N_G}\right)^{\frac{1}{3}},$$

i.e. practically for any $m_G < M_X \sim 10^{15}$ GeV. Owing to the $q\bar{q} \to G\bar{G}$ reactions G and $\bar{G}$ were in equilibrium with other particles at $T > m_G$, so that their concentration was $\sim T^3$. At $T < m_G$ the equilibrium concentration of heavy particles reduces exponentially to $\sim (m_G T)^{\frac{3}{2}} \exp[-m_G/T]$. However, this is true until the time scale of the annihilation $G\bar{G} \to q\bar{q}$ does not exceed the cosmological time scale. Afterwards the usual picture of freezing arises, giving the frozen concentration of heavy quarks

$$\nu = \frac{n_G}{n_r} \sim \frac{m_G}{\alpha m_{Pl}}, \tag{8.10}$$

where $\alpha \sim 1/50$.

When the temperature drops below T_0,

$$T_0 = \nu m_G \sim \frac{m_G^2}{\alpha m_{Pl}}, \tag{8.11}$$

the density of frozen heavy quarks $\varrho_G = m_G n_G = \nu m_G n_r$ exceeds the density of relativistic particles $\varrho_r \sim T_0 n_r$. So at

$$t_0 \sim \frac{m_{Pl}}{T_0^2} \sim \frac{\alpha^2 m_{Pl}^3}{m_G^4} \sim \alpha^2 \left(\frac{m_{Pl}}{m_G}\right)^4 t_{Pl} \tag{8.12}$$

the stage of heavy-quark dominance takes place. To survive until their dominance in the Universe, heavy quarks must have lifetime $\tau_G > t_0$.

2) In GUTs, based on gauge groups higher than SU_5, right-handed neutrinos, R-neutrinos, are predicted [69-71]. They take place in O_{10}, for example. General wisdom is to ascribe such neutrinos ν_R a very high Majorana mass M_R. If there is a Dirac mass of neutrinos m_D of the order of the Dirac masses for charged leptons, the diagonalization of the neutrino mass matrix

$$M_\nu = \begin{pmatrix} 0 & m_D \\ m_D & M_R \end{pmatrix} \tag{8.13}$$

provides a small Majorana mass $\sim m_D^2/M_R$ for left-handed neutrinos. Similar to the previous case decays of ν_R induced by exchanges of superheavy X-bosons are possible: $\nu_R \to \nu_L q\bar{q}$, etc. The lifetime of ν_R relative to these decays may be estimated by formula (8.9) (*).

If there is no other interaction of ν_R except that induced by X-bosons (it may be not the case, see footnote) soon after $T \sim M_X$, the time scale of the processes $\nu_R \nu_L \to qq$, $\ell\ell$ or $\nu_R \bar{\nu}_R \to \nu_L \bar{\nu}_L$, etc.

$$\tau \sim (n_R \sigma v)^{-1} \sim T^3 \frac{g^4}{M_X^4} T^2 N_{ch}$$

exceeds the cosmological time scale and ν_R decouple from other particles. R-neutrinos cease to interact with other particles. However, in the course of the successive adiabatic expansion their equilibrium distribution is retained until $T \sim m_R$.

They retain their equilibrium density $\sim T^3$. So the frozen concentration of ν_R is $\nu = n_{\nu_R}/n_r \sim 1/\varkappa$, where $\varkappa$ is the number of species of relativistic particles which are in equilibrium in the decoupling period of R-neutrinos. When the

(*) If m_D in (8.13) is induced by the Higgs mechanism of the Weinberg-Salam model, *i.e.* neutrino Dirac masses, as well as masses of charged leptons and quarks, are generated by interaction with the scalar field h whose self-interactions provide spontaneous breaking of the $SU_2 \times U_1$ symmetry of unified electroweak interaction, the decays $\nu_R \to \nu_L + h$ are possible. Owing to these decays, the lifetime of ν_R may be reduced substantially. The subsequent discussion contained in subsect. **8**.2.2 and is not appropriate for this case.

temperature drops below $T_0 \sim \nu m_R \sim m_R/\varkappa$ R-neutrinos dominate in the cosmological density.

3) Magnetic monopoles of mass $M \sim M_X/g \sim 10^{16}$ GeV are predicted by all the existing GUTs. These classical objects are produced in the course of GUT phase transition (see sect. 7) at $T \sim M_X$. It was shown in [53, 54] that, independently of their initial concentration (provided it is not too low), the diffusion of monopoles towards antimonopoles and their successive annihilation result in the frozen concentration

$$\nu_M = \frac{e^5 m}{m_{Pl}},$$

where e is the electromagnetic charge.

Monopoles are absolutely stable relative to decay. Annihilation with antimonopoles is the only source of their instability. As we already pointed out in sect. 7, at $T_0 \sim \nu_m m$ monopoles start to dominate in the cosmological density.

Thus all the examples show that superheavy metastable particles of mass m and relative frozen concentration $\nu = n_m/n_r$ start to dominate in the cosmological density, when the temperature drops below $T_0 = \nu m$. Dominance of nonrelativistic particles in the cosmological density means that the equation of state of the expanding Universe is no longer relativistic, $p = \varepsilon/3$. Nonrelativistic particles have a pressure negligible as compared to their energy density ε:

$$p \sim n_m m v^2$$

and $\varepsilon \sim n_m m c^2$, so $p \ll \varepsilon$ at $v \ll c$. They imply a « dustlike » equation of state $p \approx 0$.

8'2.2. Early dustlike stages and PBH formation. According to the theory of gravitational instability (see [1]) in the expanding Universe at the dustlike stage small initial density perturbations grow as

$$\left(\frac{\delta\varrho}{\varrho}\right)_M \propto \left(\frac{t}{t_M}\right)^{\frac{2}{3}} \delta_M , \tag{8.14}$$

where t_M is the moment when the perturbation with initial amplitude $(\delta\varrho/\varrho)|_{t=t_M} = \delta_M$ of given scale M enters the horizon M_h. At $t_1 \sim \delta_M^{-\frac{3}{2}} t_M$, when the value of $(\delta\varrho/\varrho)_M$ is of order 1, the nonlinear stage of the evolution of inhomogenities of scale M begins. These inhomogenities separate from the cosmological expansion and start to contract. The bulk of contracting matter maintains the structure of the inhomogenities. At the first stage flattened configurations of the pancake type may be formed. A more complicated structure is to be awaited in the course of its evolution. However, with small (but nonzero) probability a spherically symmetric homogeneous contraction may be realized. In this case, matter contracts within its gravitational radius, forming a black hole (BH). The probability of this direct BH formation is determined by the amplitude of the initial density perturbation—the smaller it is, the finest tuning of the velocity and density distributions of particles within it is needed, and thus the less probable is BH formation.

Estimates of [55, 72] have given for the probability of PBH formation

$$W_{\mathrm{PBH}} \geqslant 2\cdot 10^{-2}\,\delta_M^{13/2}\,. \tag{8.15}$$

So a very small (but nonzero) fraction of matter

$$\beta \sim \varrho_{\mathrm{PBH}}/\varrho_{\mathrm{m}} \sim 2\cdot 10^{-2}\,\delta_M'^{\frac{3}{2}}$$

goes at the dustlike stage into BHs.

The spectrum of the BH produced is determined by the spectrum of the initial inhomogeneities. However, the process of direct BH formation is effective in a fixed interval of masses—the probability of direct formation of BHs with masses smaller than the mass of the horizon at the moment t_0,

$$M_0 = m_{\mathrm{Pl}}\,t_0/t_{\mathrm{Pl}}\,, \tag{8.16}$$

is highly suppressed. This process ceases, then the dustlike stage ends. So the maximal BHs formed have the mass of the inhomogenities which started to contract at this moment, *i.e.*

$$M_{\max} = m_{\mathrm{Pl}}\,\frac{\tau}{t_{\mathrm{Pl}}}\,\delta_M^{-\frac{3}{2}}\,. \tag{8.17}$$

This mechanism provides a « tablelike » form of the PBH spectrum and may lead to, for instance, formation of PBHs evaporating before recombination only, so that no PBHs evaporating after recombination are formed. The latter possibility arises if $M_{\max} < 10^{13}$ g, *i.e.*

$$\tau\delta_M^{-\frac{3}{2}} < 10^{18}\,t_{\mathrm{Pl}}\,.$$

The mechanism provides formation of PBHs evaporating before recombination if $M_0 \leqslant 10^{13}$ g, *i.e.* if

$$t_0 < 10^{18}\,t_{\mathrm{Pl}}\,.$$

The latter is the case for (see the preceding subsection) quarks with $m_{\mathrm{G}} \geqslant 3\cdot 10^{13}$ GeV and R-neutrinos with $m_{\mathrm{R}} \geqslant 10^{11}$ GeV.

After the end of the early dustlike stage, at $t > \tau$, when the heavy particles decay and the ultrarelativistic products of their decay maintain once again the ultrarelativistic equation of state $p = \varepsilon/3$, the relative contribution of PBHs to the cosmological density grows as $(t/\tau)^{\frac{1}{2}}$, so the fraction of matter $\alpha(M)$ contained in PBHs at the moment $t_e = (M/m_{\mathrm{Pl}})^3 t_{\mathrm{Pl}}$ of their evaporation is

$$\alpha(M) = \beta(M)\left(\frac{t_e}{\tau}\right)^{\frac{1}{2}} = \beta(M)\sqrt{\frac{M^3}{m_{\mathrm{Pl}}^4\tau}}\,.$$

9. – Interaction of $\bar{\mathrm{p}}$ with ^{4}He and light-element abundances.

In sect. **4** we have shown that the study of the $\bar{\mathrm{p}}^4$He annihilation may impose a stronger restriction on the possible amount of antimatter in the early Universe than the limits that arise from the distortion of the thermal background. But the experimental investigation of the $\bar{\mathrm{p}}^4$He reaction is valuable not for only this reason. As revealed by the discussion in sect. **6-8**, from the GUT and from the physics of PBHs it follows that substantial amounts of antimatter may have appeared in the early Universe. The most important thing that one can obtain from the study of $\bar{\mathrm{p}}^4$He annihilation is the check of the possibility for annihilation to take place in the early Universe. If this is really the case, the $\bar{\mathrm{p}}^4$He annihilation must be organized in a very special way. We shall discuss this point in detail later. Now, we shall start with a brief discussion of the general features of the antiproton interactions with nuclei.

9'1. *Interactions of* $\bar{\mathrm{p}}$ *with nuclei.* – There is no experimental information on the $\bar{\mathrm{p}}^4$He interaction. For other nuclei, heavier than deuterium, the experimental data are quite scarce and incomplete. The situation with the theoretical investigations of the $\bar{\mathrm{p}}$A interaction is a little better.

In fig. 9.1 we show the typical behaviour of $\bar{\mathrm{p}}$A total and annihilation cross-sections *vs.* the energy of $\bar{\mathrm{p}}$'s calculated in the framework of Glauber approach [73]. One can see that σ_{ann} is as large as one-half of σ_{tot}.

The experimental study of the interactions of antiprotons at rest with heavier nuclei shows that, as expected, the pions from annihilation may have an appreciable chance to interact with a residual nucleus. One can draw this conclusion simply from the consideration of the average multiplicity of charged pions $\bar{n}_{\pi^\pm}$. Whereas $\bar{n}_{\pi^\pm}$ for $\mathrm{p}\bar{\mathrm{p}}$ interactions is 3.05 ± 0.04 [74], in the case of $\bar{\mathrm{p}}^{12}$C, $n_{\pi^\pm} = 2.72 \pm 0.03$ [75] and, for $\bar{\mathrm{p}}^{208}$Pb, $n_{\pi^\pm} = 2.44$ [76].

In fig. 9.2 we show the multiplicity distribution for « hadronlike » particles emitted from emulsion nuclei after the annihilation of stopped antiprotons. The experimental data are from [76]. The histogram is calculated in the ex-

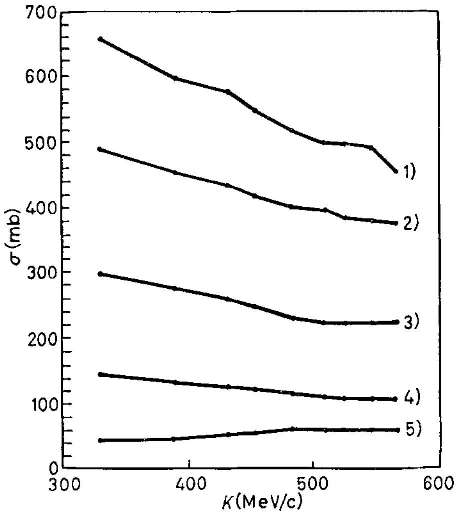

Fig. 9.1. – Momentum dependence of the $\bar{p}^3He$ cross-sections: curve 1) is for the cross-section without shadowing, *i.e.* for the sum $2\sigma_{tot}(\bar{p}p) + \sigma_{tot}(\bar{p}n)$; curve 2) is for the total cross-section; curve 3) is for the annihilation cross-section; curve 4) is for the total elastic cross-section and curve 5) is for the total cross-section of ^3He break-up processes.

tensive paper [77], in which a number of aspects of the $\bar{p}A$ interaction at rest was studied by means of a Monte Carlo simulation.

From fig. 9.2 one can see that up to 20 secondary particles may be emitted from a nucleus. (We recall that neutrons were not detected in the emulsion, so the number of secondaries must be even larger.)

More detailed information about the interaction of pions in final states is presented in table 9.I, where the probabilities of pion interactions of different orders are summarized [77].

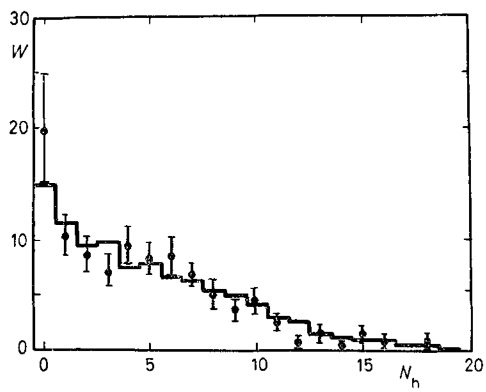

Fig. 9.2. – The multiplicity distribution (in percent) for hadronlike particles emitted by the nuclei of the photoemulsion after the annihilation in it of stopped antiprotons.

TABLE 9.I. – *Probabilities of pion interactions of different orders for the absorption of* $\bar{p}$'s *by nuclei* ^{70}Ga *and* ^{208}Pb (*from paper* [77]). ν is the total number of pions interacting with a nucleus.

ν	0	1	2	3	4	5	6	7	$\bar{\nu}$
^{70}Ga	17	24.3	27	19.5	9.8	2.6	0.8	—	1.91
^{208}Pb	13.2	20	28.7	22	10.7	4.1	1.0	0.2	2.15

It is seen that on the average a pion undergoes two interactions with a nucleus but with a probability of a few percents it may interact $5 \div 6$ times.

Therefore, the final interaction of pions from annihilation is very important and it may change the whole picture of the $\bar{p}A$ interaction substantially. One may expect the outputs of nucleus fragments due to these pion interactions to differ from those obtained on the basis of a naive one-nucleon mechanism of the $\bar{p}A$ interaction.

9·2. *The deuterium puzzle.* - The so-called deuterium puzzle is the bridge which links together the properties of $\bar{p}^4$He annihilation and its astrophysical consequences.

In fig. 9.3 are presented the concentrations of different light elements calculated within the framework of the standard big-bang cosmology [78, 79]. One can see that the deuterium abundance depends strongly on the density of baryons ϱ_B in the Universe or on the parameter $\Omega_B = \varrho_B/\varrho_c$, where ϱ_c is the critical density. The parameter Ω_B is not well known (the famous paradox of the « hidden mass »). From one type of observations, based on the measurements of the luminosity of stellar objects, Ω_B is less than 0.1. From another type of observations, based on the measurements of the relative velocities of the galaxies which depend on the gravitational masses of the galaxies, the value of Ω_B is of the order of $0.2 \div 0.7$. The observed abundance of deuterium $X_d = (2.5 \pm 1.5) \cdot 10^{-5}$ [36] is consistent with $\Omega_B < 0.1$, but in the case of $\Omega_B \geqslant 0.2 \div 0.7$ the concentration of deuterium must be of the order of $X_d \sim 10^{-8} \div 10^{-6}$. There exist weighty arguments to assume that Ω_{tot} is really $\sim 0.2 \div 0.7$. Therefore, if $\Omega_{tot} \equiv \Omega_B$, additional sources of deuterium, besides big-bang nucleosynthesis, are needed to explain the observed abundance of deuterium. A number of attempts has been made to solve this problem, but all in vain.

For example, GOLDBERG and CHECHETKIN [80] considered the formation of deuterium due to spallation of ^{4}He by protons and in $\alpha + \alpha$ reactions. It must be mentioned that to spallate ^{4}He a proton is required to have an energy $\geqslant 28$ MeV. From the point of view of particle physics that is quite a small energy, easily obtained in accelerators. But there are only few processes in the space which may provide a substantial flux of tens-of-MeV protons. In [80] the processes in the envelope of supernovae were considered. It was shown that

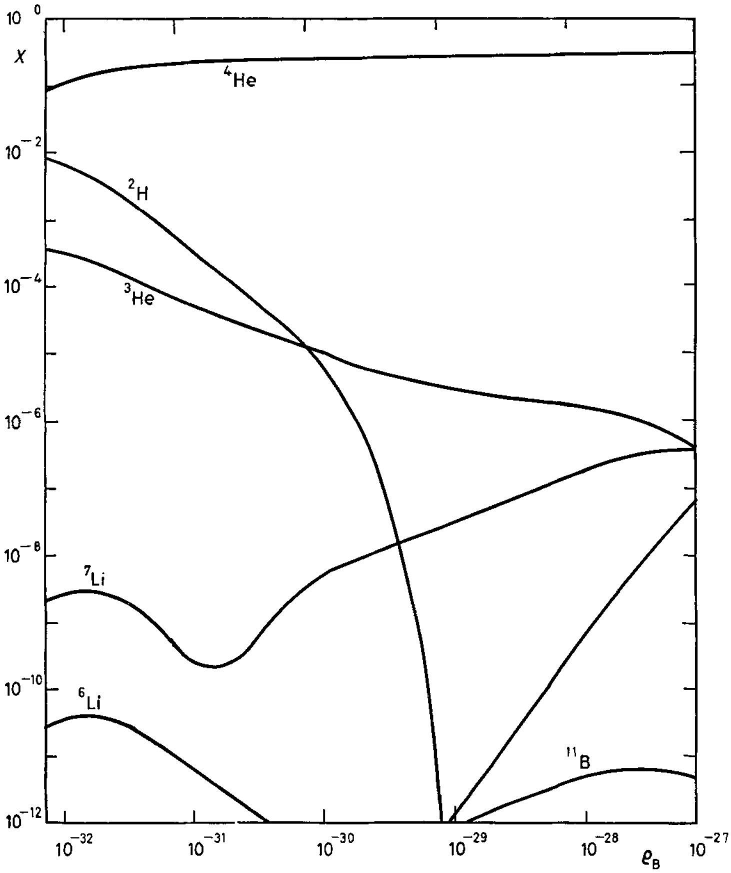

Fig. 9.3. – Predictions of the primordial abundances of d, ^{3}He, ^{4}He, ^{6}Li, ^{7}Li, ^{11}B for different baryon densities ϱ_b. From the results of calculations [78, 79].

the assumption that additional deuterium is formed in p^{4}He spallation and $\alpha+\alpha$ reactions implies large amounts of ^{3}He, ^{6}Li and ^{7}Li being formed, too. For example, the calculated ratio of X_d/X_{Li} is two or three orders less than the observed one.

The deuterium may be formed due to $\bar{p}^4$He annihilation, too. Even antiprotons at rest may create deuterium in $\bar{p}^4$He annihilation. From the previous consideration it is clear that, not only the pure output of deuterium is significant, but also the relationship between the outputs of other light elements is extremely important, too.

9'3. *May* $\bar{p}^4$He *data prove the existence of late annihilation in the Universe?* – In this subsection we discuss the question if the $\bar{p}^4$He data can prove or disprove the possibility of late annihilation in the Universe. It is clear that the observed abundances of d and ^{3}He put some limits on the outputs of these nuclei in

$\bar{p}^4$He annihilation. For example, if we suggest that all deuterium in the Universe was formed in $\bar{p}^4$He annihilation and it turns out to be that the output of ^{3}He in this annihilation is 100 times greater than the output of d, that definitely rules out the first suggestion. Remember that the observed mass fractions of d and ^{3}He are the following:

$$X_d = (2.5 \pm 1.5)\cdot 10^{-5}\,, \qquad X_{^3\mathrm{He}} = (4.2 \pm 2.8)\cdot 10^{-5}\,. \tag{9.3.1}$$

Therefore, the outputs of d and ^{3}He in $\bar{p}^4$He annihilation cannot differ substantially from one another. Otherwise, we should see a too large amount of ^{3}He, in our example 100 times greater than the observed abundance of ^{3}He. So, if one observes a large difference between the outputs of d and ^{3}He in $\bar{p}^4$He annihilation, one can immediately conclude, from this fact, that annihilation hardly took place in the early Universe.

Let us try to obtain more or less rigorous limits on the ratio of f_d^{eff} and $f_{^3\mathrm{He}}^{\mathrm{eff}}$, which are the outputs of d and ^{3}He in $\bar{p}^4$He annihilation. Keep in mind that in the hot early Universe (see sect. **4**)

$$f_d^{\mathrm{eff}} = f_n + f_d \qquad\qquad \text{at } t \leqslant t_d$$

and then at $t > t_d$ (free neutrons do not succeed in forming deuterium in the $n+p \to d+\gamma$ reaction)

$$f_d^{\mathrm{eff}} = f_d\,.$$

Let us suppose that a part of the observed deuterium (as well as ^{3}He) was formed in $\bar{p}^4$He annihilation X_d^{ann} and the other part $X_d(\varrho_B)$ was formed due to the big-bang nucleosynthesis. The symbol ϱ_B in the parentheses signifies that the amount of « big-bang » deuterium depends on the density of baryons in the Universe. When ϱ_B changes from $\sim 5\cdot 10^{-31}$ to $5\cdot 10^{-30}$ g/cm^3, $X_d(\varrho_B)$ varies from $\sim 5\cdot 10^{-5}$ to $\sim 10^{-7} \div 10^{-8}$ (see fig. 9.3). At a first glance, the suggestion that all the observed deuterium X_d^{obs} was formed due to annihilation implies that we must deal with the case of large ϱ_B, $\Omega_B \geqslant 0.1$, *i.e.* $\varrho_B \geqslant 0.1\,\varrho_c$ at which $X_d(\varrho_B)$ is rather small. So

$$X_d^{\mathrm{obs}} = X_d^{\mathrm{ann}} + X_d(\varrho_B)\,,$$

where $X_d^{\mathrm{ann}} \sim 10^{-5}$ and $X_d(\varrho_B) \sim 10^{-7} \div 10^{-8}$.

Actually, this is not necessary. The generous errors in X_d^{obs} and $X_{^3\mathrm{He}}^{\mathrm{obs}}$ (see (9.3.1)) allow as large an amount of annihilated deuterium X_d^{ann} as the « big-bang » one:

$$X_d(\varrho_B) \approx X_d^{\mathrm{ann}} \approx X_d^{\mathrm{obs}}\,.$$

It must be noted that the quoted errors in (9.3.1) are not the errors of some definite measurement but the average over the whole set of observations. So these errors to a greater extent reflect not the precision of a certain measurement but the deviation of X_d in different regions of space (see appendix, table A.I).

The situation with $X_{^3\mathrm{He}}$ is the same (if not worse). Nevertheless, we may draw some important conclusions, in spite of the uncertainties in $X^{\mathrm{obs}}_{^3\mathrm{He}}$ and $X^{\mathrm{obs}}_{\mathrm{d}}$. Let us assume that $\bar{\mathrm{p}}^4$He annihilation in the early Universe provides $X^{\mathrm{ann}}_{\mathrm{d}} \sim X^{\mathrm{obs}}_{\mathrm{d}}$ and

$$0 \leqslant X^{\mathrm{ann}}_{^3\mathrm{He}} \leqslant 10 X^{\mathrm{obs}}_{^3\mathrm{He}} \,. \tag{9.3.2}$$

(The upper limit approximately corresponds to an interval of three « standard deviations » for $X^{\mathrm{obs}}_{^3\mathrm{He}}$. The lower limit indicates that the amount of ^{3}He from $\bar{\mathrm{p}}^4$He annihilation is negligible, and all the ^{3}He is due to the big-bang nucleosynthesis.)

Then the ratio between the outputs of d and ^{3}He in $\bar{\mathrm{p}}^4$He should be the same as that between $X^{\mathrm{ann}}_{\mathrm{d}}$ and $X^{\mathrm{ann}}_{^3\mathrm{He}}$:

$$\frac{f^{\mathrm{eff}}_{\mathrm{d}}}{f_{^3\mathrm{He}}} = \frac{3}{2}\,\frac{X^{\mathrm{ann}}_{\mathrm{d}}}{X^{\mathrm{ann}}_{^3\mathrm{He}}} \,. \tag{9.3.3}$$

Taking into account (9.3.2), one can obtain

$$0.089 < \frac{f^{\mathrm{eff}}_{\mathrm{d}}}{f_{^3\mathrm{He}}} < \infty \,. \tag{9.3.4}$$

Therefore, if in the experiments on $\bar{\mathrm{p}}^4$He annihilation it will be found that the output of ^{3}He is less than the effective output of d, one can conclude that this result does not contradict the possibility of late annihilation (at $10^3\,\mathrm{s} < t < 10^{13}\,\mathrm{s}$) in the Universe. If, otherwise, it will be shown that the output of ^{3}He is greater than the deuterium one, that rules out the possibility of late annihilation. The latter statement needs some comments. Strictly speaking, the condition (9.3.2) that $X^{\mathrm{ann}}_{\mathrm{d}} \sim X^{\mathrm{obs}}_{\mathrm{d}}$ is not necessary. One may speculate that not all the deuterium was formed in annihilation, but only a certain part of it. So it will be found that, if $f_{^3\mathrm{He}} > f_{\mathrm{d}}$, we may put some restrictions on the part of deuterium being formed in $\bar{\mathrm{p}}^4$He annihilation and, ultimately, we put the restrictions on the amount of antimatter in the early Universe. It must be stressed that this restriction will be stronger than that from eq. (4.7) in sect. **4**.

9'4. *Post big-bang production of light elements.* – A comparison of observational data (see the appendix) on the abundances of light ($A < 12$) elements with the results of numerical calculations [78, 79] of primordial nucleosynthesis leads to the conclusion that some additional (post big-bang) sources of light elements must exist. These sources are usually related to nonstationary processes in the Galaxy, such as supernova explosions, cosmic-ray interactions, etc. Let us first consider the possibility of deuterium production by Galaxy sources. We have mentioned in the preceding subsection the idea of [80] and [81] on the deuterium formation in the shock waves in supernova explosions as a result of ^{4}He spallation within the front of the shock wave. Subsequent calculations [80] of the complete set of nuclear reactions in this process have led to the conclusion that the amount of ^{6}Li produced as a by-product in this process is so large that it is in contradiction with its observed abundance.

In paper [82] the process of ^{4}He spallation in the thick disk formed at the accretion of the matter on the black hole was suggested. The other possibility of deuterium formation is connected with neutron stars. Free neutrons may form deuterium, whether the neutron star is disrupted owing to tidal effects, or the neutron-rich matter formed in its interior is thrown out on its surface [83]. However, there was no quantitative analysis of these processes of deuterium formation.

As was mentioned above (sect. **2**), one of the essential features of the chemical evolution of light elements is the burning of deuterium into ^{3}He in the stars. The amount of burned deuterium is determined by the conditions inside the star. There is no quantitative answer to this question, but we may estimate the upper limit of the primordial deuterium abundance with the use of the total concentration $\mathrm{D}+{}^3\mathrm{He}$ (as was done in ref. [1]) for the estimation of the deuterium abundance in the protosolar cloud.

Another essential mechanism of light-element formation is connected with the cosmic-ray interactions. This mechanism implies the spallation of the ^{12}C and ^{16}O nuclei (the most abundant ones among the nuclides produced in stars) by cosmic rays. For an extensive review on this subject see [84]. Li, B, Be, ... seem to be produced by this mechanism. The main problem in the quantitative estimates of the outputs of these elements is the question on the intensity of cosmic rays and on the time variation of this intensity. Based on the intensity of cosmic rays observed in the neighbourhood of the solar system the observed abundances of ^{10}Be and ^{11}B can hardly be explained in the framework of the considered mechanism.

The possibility of ^{11}B formation in the envelopes of collapsing stars due to the interaction of the ν_e flux with ^{12}C nuclei was considered in [85]. $\bar{\mathrm{p}}$ annihilation with ^{12}C and ^{16}O nuclei, discussed in the next subsection, may be treated as an interesting mechanism of Li, B, Be formation in the interstellar medium.

9'5. *Annihilation at the Galaxy stage and formation of light elements.* – The sources of late annihilation predicted by GUTs (see sect. **8**) may work effectively at the post recombination stage (at $t \geqslant 10^{13}$ s) either. PBHs with a mass of $10^{13}\,\mathrm{g} < M < 10^{15}\,\mathrm{g}$ evaporate in this period. Antimatter domains of scale $N_B \geqslant 10^{67} r_B \bar{r}$ annihilate in this period. So a slight change in the parameters of GUTs, determining the masses and the lifetimes of superheavy metastable particles (subsect. **8**˙1), and thus the spectrum of PBHs, or the scale of antimatter domains (subsect. **8**˙2), may result in the prediction of the appearance of a substantial amount of antimatter (or antinucleons) at the post recombination stage.

To discuss the tests for the possible presence of these sources of antimatter, let us consider the relationship of $\bar{p}$ annihilation and light-element abundances at the post recombination stage (*i.e.* at $Z < 10^3$, or $t > 10^{13}$ s). The data on the γ-background (see sect. **3**) put limits on the possible fraction f of annihilated antinucleons. These limits are very strong if annihilation takes place at present ($f < 10^{-15} \div 10^{-10}$ in the Galaxy, $f < 10^{-5}$ in the intergalactic medium, see **3**˙2.2), but they become weaker with the increase of Z. If annihilation had finished, say, at $Z \sim 1 \div 2$, there would have been no contradiction with the γ-background if $f \sim 10^{-3}$. According to the modern theory [1] of the evolution of the Universe, the stars are formed at $Z \leqslant 5 \div 10$. As a result of stellar evolution and successive stellar disruption heavy elements are produced ($A \geqslant 12$). Antiproton interactions with such elements may provide formation of rare isotopes with relatively low abundances. The rare elements with $A < 12$ are of special interest here, since the observed abundances of Li, B, Be might have been thus connected with the processes

$$\bar{p} + {}^{12}C \to \left\{ \begin{array}{l} {}^{7}Li + \text{anything}, \\ {}^{6}Li + \text{anything}, \\ {}^{9}Be + \text{anything}, \\ {}^{10}B + \text{anything}, \\ {}^{11}B + \text{anything}, \end{array} \right. \tag{9.5.1}$$

$$\bar{p} + {}^{16}O \to \left\{ \begin{array}{l} {}^{6}Li + \text{anything}, \\ {}^{7}Li + \text{anything}, \\ {}^{9}Be + \text{anything}, \\ {}^{10}B + \text{anything}, \\ {}^{11}B + \text{anything}. \end{array} \right. \tag{9.5.2}$$

Let us write down the kinetic equation for these processes:

$$\frac{1}{A}\frac{dX_A}{dt} = \frac{X_{C,O} f(Z)}{A_{C,O}} n_B (1 + Z)^3 \langle \sigma v \rangle_{C,O}, \tag{9.5.3}$$

where X_A, X_C or X_O are the mass concentrations of the isotopes A, carbon and oxygen, respectively, n_B is the baryon density, $f(Z)$ is the fraction of annihilated $\bar{p}$'s and $\langle\sigma v\rangle_{C,O}$ is the mean rate of reactions (9.5.1) and (9.5.2). Recalling eq. (3.2.1) for the γ-flux F_γ produced in the annihilation of the fraction f of $\bar{p}$'s with the matter and taking into account the red-shift Z, one obtains

$$F_\gamma = k_\gamma f(Z) n_B^2 (1+Z)^3 \sigma_0 c t_{ann}(Z)\,, \tag{9.5.4}$$

where $k_\gamma \sim 5$ is the mean number of produced γ's per annihilation,

$$n_B = \Omega_B \cdot 10^{-6}/\mathrm{cm}^3$$

is the modern baryon number density, $t_{ann}(Z)$ is the time scale of annihilation at the red-shift Z, and the cross-section of annihilation of slow $\bar{p}$'s is taken as $\sigma_{ann} = \sigma_0 c/v$, so that $\sigma_{ann} v = \sigma_0 c$. Combining eqs. (9.5.3) and (9.5.4), one obtains the following expression for the increase of the abundance ΔX_A of the element A as a function of Z:

$$\Delta X_A = 4\,\frac{A}{A_{C,O}}\,\frac{\sigma v\left(\bar{p}+{C \atop O} \to A+\text{anything}\right)}{\sigma_0 c}\,\frac{F_\gamma(100\ \mathrm{MeV}/(1+Z))}{n_\gamma c} \tag{9.5.5}$$

or for $\sigma v\left(\bar{p}+{C \atop O} \to A+\text{anything}\right) = \sigma^{tot}\left(\bar{p}+{C \atop O}\right) v f_A \sim A_{C,O}^{\frac{2}{3}}\,\sigma_0 c f_A$

$$\Delta X_A \sim f_A A^{\frac{2}{3}} \cdot 4 \cdot 10^{-10} X_{C,O}\left[\frac{F_\gamma(100\ \mathrm{MeV}/(1+Z))}{10^{-5}/\mathrm{cm}^2\,\mathrm{s}}\right]. \tag{9.5.6}$$

If annihilation took place at $Z=5$, one obtains (from the existing limit $F_\gamma(20\ \mathrm{MeV}) \leqslant 4\cdot 10^{-4}/\mathrm{cm}^2\,\mathrm{s}$) for $X_{C,O} \sim 10^{-2}$

$$\Delta X_A \leqslant f_A A_{C,O}^{\frac{2}{3}} \cdot 2 \cdot 10^{-10}\,. \tag{9.5.7}$$

One obtains from (9.5.7) that, in the case of Be, $\Delta X_{Be} \leqslant 4.5 f_{Be} X_{Be}^{obs}$. So, if $f_{Be} \geqslant \frac{1}{5}$, the observed abundance of beryllium may be explained by the $\bar{p}^{12}C$ or $\bar{p}^{16}O$ annihilation. The obtained relationship demonstrates the dependence of the possible output of the element A and $\bar{p}C(O)$ annihilation on the experimentally measurable quantity

$$f_A = \frac{\sigma v\left(\bar{p}+{C \atop O} \to A+\text{anything}\right)}{\sigma v_{\bar{p}p}}\,. \tag{9.5.8}$$

The measurement of the cross-sections of reactions (9.5.1) and (9.5.2) provides a definite check of the hypothesis on the $\bar{p}+C$ or $\bar{p}+O$ origin of the Li, Be, B abundances. To prove (or exclude) the possibility of the annihilation nature of light elements the outputs of all the nuclides in reactions (9.5.1) and (9.5.2) are of importance. Similarly to the discussions of the preceding subsect. **9**.3 (and **9**.4) on the $\bar{p}^4$He nature of the observed deuterium, the $\bar{p}^{12}$C or $\bar{p}^{16}$O nature of the observed Be may be checked both by the magnitude of its output and by the correlation of its output with the output of other nuclides. Here we do not consider this problem in detail, our aim is to demonstrate that the measurement of the outputs of less abundant nuclides in the $\bar{p}$ interactions with relatively more abundant nuclides may serve as a useful tool for the check of the recent presence of the sources of antiprotons (antimatter) in the Universe predicted by GUTs.

Note, however, that « natural » sources of antiprotons, such as interactions in the neighbourhood of pulsars or in the accretion disks around the massive black holes in the centres of galaxies, arise at the galaxy formation stage. This fact complicates the interpretation of the mechanism of $\bar{p}$ production, leaving intact the necessity of an experimental check of the « annihilation » mechanism of rare-element production in the not so far past.

10. – What new information can be obtained from the study of $\bar{p}$A annihilation.

In this section we summarize all results which partially were obtained in the previous sections of our review.

The most promising is the study of the antiproton annihilation with ^{4}He. By measuring the effective output of ^{3}He ($f^{\text{eff}}_{^3\text{He}}$) in $\bar{p}^4$He annihilation one may obtain a restriction on the fraction f of antimatter in the early Universe at $10^3\,\text{s} \leqslant t \leqslant 10^{13}\,\text{s}$. The measurements of the effective output of D in $\bar{p}^4$He annihilation may provide more detailed information. One can obtain a restriction on the fraction of antimatter in the Universe in two periods, the first one from 10^3 s up to approximately 1 month from the beginning, and the second from 1 month up to 10^{13} s. As we mentioned in sect. **4**, these restrictions are, at least, four orders of magnitude stronger than that which is obtained from γ-radiation observations. Besides that, by comparison between $f^{\text{eff}}_{\text{d}}$ and $f^{\text{eff}}_{^3\text{He}}$, one can draw the conclusion about the very possibility for the late annihilation in the Universe.

As we noted in sect. **9**, if

$$0.89 \leqslant f^{\text{eff}}_{\text{d}}/f^{\text{eff}}_{^3\text{He}} \leqslant \infty\,, \tag{10.1}$$

the annihilation in the Universe does not contradict the observed abundances of D and ^{3}He. Otherwise, even more stringent restrictions on f could be derived.

The estimation of f based on the experimental information on $\bar{p}^4$He annihilation may lead to a number of interesting consequences.

10·1. *Limits on the relative contributions of the primordial black holes* $\alpha(M)$ *and* $\beta(M)$. From eq. (6.14) one can obtain that

$$\alpha(M) = \frac{f\Omega_B(t_{ev}/1\,\mathrm{s})^{\frac{1}{6}}}{25f_{\bar{B}}}, \tag{10.2}$$

where $\alpha(M) = \varrho_{PBH}/\varrho_{tot}$, $f_{\bar{B}} \sim 0.1$ is the fraction of antibaryons in the radiation spectrum of PBHs. So, if we assume that $f \sim 10^{-3}$ (see sect. **4**), $t_{ev} \sim 10^6$ s and $\Omega_B \sim 0.1$, then

$$\alpha(M) \leqslant 0.4 \cdot 10^{-3}. \tag{10.3}$$

The corresponding restriction deduced from the distortions of the thermal-background radiation is two orders of magnitude weaker (see sect. **3**). Restrictions on the value of $\alpha(M)$ based on the $\bar{p}^4$He data can give information about the homogeneity of the early Universe at $t \geqslant 10^{-28}$ s. These restrictions may put stringent limits on the inhomogeneities appearing in the course of GUT phase transition.

Indeed, owing to the enhancement of the relative contribution of PBHs to the cosmological density at the RD stage and at the stages of relativistic-particle dominance (at $t < 1$ s), the restriction (10.3) on the value of $\alpha(M = 10^{11}\,\mathrm{g})$ converts into the restriction on the probability W_{PBH} of formation of such PBHs at $t \sim 10^{-27}$ s:

$$W_{PBH} \sim \beta(M) = \frac{m_{Pl}}{M}\alpha(M) \leqslant 10^{-19}, \tag{10.4}$$

where $m_{Pl} = 10^{-5}$ g is the Planck mass.

10·2. *Limits on the* GUT-*predicted superheavy particles.* – Rigorous restrictions on the PBH formation probability result in restrictions on the masses and lifetimes of GUT-predicted superheavy metastable particles. The existence of these particles leads (see sect. **8**) to the possibility of PBH formation from small density perturbations at the stage of the dominance of such particles in the early Universe. At the initial inhomogeneity $\delta\varrho/\varrho \equiv \delta \sim 10^{-2}$, particles with mass $m \sim 10^{13}$ GeV, dominating in the Universe at $t \geqslant 10^{-27}$ s, form PBHs of mass $M \sim 10^{11}$ g with the probability

$$W_{PBH}^{(p)} \geqslant 2 \cdot 10^{-2}\delta^{\frac{13}{2}} \sim 2 \cdot 10^{-15}. \tag{10.5}$$

In the case of superheavy-metastable-particle dominance one has (see sect. **8**) a relationship between the probability $W^{(\mathrm{p})}_{\mathrm{PBH}}$ and the value of $\alpha(M)$:

(10.6) $$\alpha(M) = \left(\frac{M}{m_{\mathrm{Pl}}}\right)^{\frac{3}{2}} \left(\frac{t_{\mathrm{Pl}}}{\tau}\right)^{\frac{1}{2}} W^{(\mathrm{p})}_{\mathrm{PBH}},$$

where τ is the lifetime of the particles.

So in the considered case limits on the value of $\alpha(M)$ put restrictions on the lifetime of the particles $\tau > 10^{-19}$ s. However, a typical lifetime for GUT-predicted particles of mass $m \sim 10^{13}$ GeV is given by (8.9), so that $\tau \sim 10^{-22}$ s.

So restriction (10.3) practically excludes either the existence of superheavy metastable particles with mass $\sim 10^{13}$ GeV, or (since the probability (10.5) depends strongly on δ) the existence of perturbations with $\delta \geqslant 10^{-2}$ in the early Universe.

Note that this result is obtained owing to restriction (10.3) on the value of $\alpha(M)$. For the restriction given from the observations of the thermal-background spectrum $\alpha(M) \leqslant 10^{-1}$, so that the respective lower limit for τ is $\tau > 10^{-23}$ s—the value $\tau \sim 10^{-22}$ s is possible and no contradictions arise. We see that the two-order-of-magnitude enhancement of the restriction on $\alpha(M)$, based on $\bar{\mathrm{p}}^4$He data, results in the new possibilities of tests of GUTs, impossible by other means.

10'3. *Limits on the annihilation of antimatter domains.* – In the considered stage $10^3\,\mathrm{s} < t < 10^{12}$ s antimatter domains of scale $N_{\bar{\mathrm{B}}} \leqslant 1.6 \cdot 10^{67} \bar{r} r_{\mathrm{B}}$ (see sect. **6**) dissipate (where r_{B} is $n_{\mathrm{B}}/n_{\gamma}$, *i.e.* the baryon-to-photon ratio). The $\bar{\mathrm{p}}^4$He data will provide a limit on the possible relative amount of antimatter f contained in such domains:

(10.7) $$f \leqslant 10^{-4}.$$

There are practically no restriction on f contained in such domains from the observations of the thermal-background spectrum

(10.8) $$f \leqslant 1 \quad \text{at } Z > 10^4 (t < 10^{10}\,\mathrm{s}).$$

10'4. *Limits on the parameters of* GUT *phase transitions and mechanisms of CP violation.* – Within the frame of GUTs antimatter domains may appear in the course of GUT phase transition, if there is soft *CP* violation (sect. **8**). The scale of a domain is determined by the duration of the phase transition (or the succession of transitions), and the value of f is determined by the relationship between the phases of hard, φ_{h}, and soft, φ_{s}, *CP* violation

(10.9) $$f = \frac{\varphi_{\mathrm{s}} - \varphi_{\mathrm{h}}}{\varphi_{\mathrm{s}} + \varphi_{\mathrm{h}}}.$$

If a domain of scale

$$(10^{45} \div 10^{62}) f r_{\mathrm{B}} \leqslant N_{\bar{\mathrm{B}}} \leqslant 1.6 \cdot 10^{67} f r_{\mathrm{B}} \tag{10.10}$$

is formed in the GUT transitions, restriction (10.9) provides tight bounds on the phases φ_s and φ_h—they must be very close to each other:

$$\left| \frac{\varphi_h - \varphi_s}{\varphi_h + \varphi_s} \right| \leqslant 10^{-4} . \tag{10.11}$$

So definite restrictions on the GUT mechanism of *CP* violation may be provided by the $\bar{p}^4$He data. There is almost no restriction on such mechanisms from the observations of the thermal-background spectrum.

10·5. *New mechanism of deuterium formation.* – The sources of late $\bar{p}$ annihilation predicted by GUTs are characterized by the limited time interval of their switching on. PBHs in a definite mass interval are formed. Domains of definite scale are predicted. The parameters of the sources of antimatter are determined by the parameters of GUTs, so that for a certain choice of these parameters antiprotons (antimatter) may appear at the RD stage only, and the deuterium output in the $\bar{p}^4$He annihilation may be just the one being observed.

The outputs of all the other products (n, ^{3}He, T) are to be measured in the $\bar{p}^4$He annihilation to check this mechanism. A special, very restricted, choice of GUT cosmology is needed for this mechanism.

Note that the presence of sources of late annihilation at the galaxy formation stage may be checked by measurements of the outputs of relatively rare elements in the $\bar{p}$ interactions with relatively abundant nuclides (C, O, Fe, etc.).

In conclusion the study of $\bar{p}$A interactions may provide *a*) limits on the possible amount of antimatter in the early Universe, *b*) limits on the probability of PBH formation and on the properties of superheavy metastable particles, *c*) restrictions on the GUT parameters determining the properties of domains of antimatter, *d*) new sources of rare elements.

* * *

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Appendix

Light-element abundances.

The observational data and the calculations of the big bang nucleosynthesis. The observational data will be presented for different elements in the order of growth of their atomic number. All the abundances taken from ref. [1, 36, 83, 84, 86, 87] will be given relative to the hydrogen one.

Deuterium. The deuterium abundance in water is $D/H = 1.6 \cdot 10^{-4}$. A similar concentration was observed in the meteorites. Estimates of the abundance of deuterium in Jupiter's atmosphere give $D/H = (2.9 \div 7.5) \cdot 10^{-5}$ from the observations of CH_3D molecules and $D/H = (2.1 \pm 0.4) \cdot 10^{-5}$ from the observations of the HD molecular content. The observations of the spectrum of the solar photosphere and corona give an upper limit on the deuterium abundance of $D/H \leqslant 4 \cdot 10^{-6}$. The absence of deuterium on the solar surface may be interpreted as an effect of deuterium burning. The observed concentration of D in water is to be corrected by the effect of its enrichment owing to molecular-exchange reactions.

Observations of deuterium in the interstellar medium indicate a rather large variation of its concentration. So, for instance, the deuterium abundance in the Orion nebula is $D/H \sim 6 \cdot 10^{-3}$ (in DCH molecules), while observations of Lyman lines in the interstellar gas give $D/H = (1.4 \pm 0.2) \cdot 10^{-5}$. The variation in the measured values of the deuterium abundance does not mean a bad accuracy of the measurements, but reflects the real inhomogeneity in the distribution of D. This inhomogeneity may be induced by some processes, leading to enrichment of deuterium.

The average magnitude of the deuterium abundance at the moment of the solar-system formation is accepted as $(2.5 \pm 1.5) \cdot 10^{-5}$ [36]. Though there is some uncertainty in the account for deuterium burning in stars (the fraction of matter participating in the stellar evolution is not well established) and in the account for the processes of deuterium production in the astrophysical objects (for example, in the inequilibrium layer of neutron stars), we may assume that the deuterium abundance has not changed drastically since the period of solar-system formation.

^{3}He. The abundance of helium-3 in the interstellar medium is determined by observations of the superfine splitting of the line 3.46 cm. This method has given the following upper limits on the ^{3}He abundance: $^3\text{He}/\text{H} < 5 \cdot 10^{-5}$. Direct observations of the ^{3}He abundance in the solar wind give $^3\text{He}/^4\text{He} = 4 \cdot 10^{-4}$. The analysis of the concentration of isotopes of some other elements on the solar surface leads to the conclusion that the value of the ratio $^3\text{He}/^4\text{He}$ cannot be essentially reduced by mixing or thermonuclear processes. Thus the observed abundance of helium-3 cannot exceed the sum of deuterium and helium-3 abundances in the protosolar gas. We may conclude that the abundance of helium-3 is $(4.2 \pm 2.8) \cdot 10^{-5}$ [36].

^{4}He. There is a large amount of estimates of ^{4}He abundances, based on the observations of excited He in galaxies and stars, as well as on the analysis of the

^{4}He abundance in the stellar material. A detailed discussion of these data is beyond to scope of the present review. As a conservative estimate of the ^{4}He abundance, we will take $^4\text{He/H} = 0.10 \pm 0.02$ [88,89]. In our opinion stringent limits on the primordial ^{4}He abundance deduced from the observed ^{4}He abundance do not seem reliable, due to ambiguities in the interpretation of observations.

^{6}Li *and* ^{7}Li. These elements were observed in different kinds of astronomical objects: in meteorites, stars, in the interstellar gas. In the latter case, the absorption line 6708 Å was observed. We may conclude from these observations that the protosolar Li/H abundance is $10^{-9\pm3}$ with isotopic composition $^7\text{Li}/^6\text{Li} \sim 12$. There were no substantial changes of these values since the period of solar-system formation. The observed variation of the Li abundance in various stars (cf. $\text{Li/H} = 10^{-7}$ in red giants, and by hundred times smaller in the Sun) is related to the concrete thermonuclear reactions and physical processes in these stars.

^{9}Be *and* ^{10,11}B. According to the modern cosmological ideas these elements could not be formed in any sizable amount in the cosmological nucleosynthesis. The cosmic-ray–induced spallation reactions on heavier nuclei (^{12}C, ^{16}O) are considered usually as the main mechanism of their formation. The observational data on the abundance of these elements are given in table A.I.

TABLE A.I. – *Abundances of light elements.*

Relative concentration	Interstellar gas	Solar surface	Stars	Meteorites
^{2}D/H	$< 4\cdot 10^{-4}$ $> 3\cdot 10^{-5}$	$< 4\cdot 10^{-6}$	$< 6\cdot 10^{-4}$	$< 8\cdot 10^{-5}$ $> 1\cdot 10^{-5}$
Accepted value: $1.4\cdot 10^{-5}$				
^{3}He/H	$< 5\cdot 10^{-5}$	$< 4\cdot 10^{-5}$		$1.5^{+1.5}_{-0.7}\cdot 10^{-5}$
Accepted value: $(2 \pm 1)\cdot 10^{-5}$				
^{4}He/H	$0.11 + 0.03$	~ 0.10	~ 0.10	
Accepted value: 0.10 ± 0.02				
Li/H	$3\cdot 10^{-10}$	10^{-11}	$10^{-10} \div 10^{-7}$	$1.5\cdot 10^{-9}$
Accepted value: 10^{-9}				
Be/H	$< 7\cdot 10^{-11}$	10^{-11}	$4\cdot 10^{-11} \div 10^{-12}$	$2\cdot 10^{-11}$
Accepted value: $2\cdot 10^{-11}$				
B/H	$< 2\cdot 10^{-9}$	$< 6\cdot 10^{-10}$		$2\cdot 10^{-10} \div 6\cdot 10^{-9}$
Accepted value: $5\cdot 10^{-9}$				

In table A.I the observational data on the abundances of all the light elements are presented. The data are taken mainly from review [86].

The standard model of big-bang nucleosynthesis. We take as a standard the baryon-asymmetric model with the ratio of baryon n_B to photon n_γ number densities $n_B/n_\gamma \sim 10^{-9\pm1}$. The most complete calculations of the primordial nuclesynthesis were performed in [78, 79] (see also [90-93]). We briefly discuss here the main assumptions, parameters and results of these calculations.

Assumptions:

1) The expansion of the Universe is considered within the framework of the metric theory of gravity.

2) At the stage of nucleosynthesis there is no strong anisotropy and inhomogeneity of the Universe.

3) There has been a high-temperature stage $T \geqslant 10^{11}$ K in the evolution of the Universe. At this stage, the equilibrium between nucleons, neutrinos, radiation and electron-positron pairs was maintained. In the course of the successive expansion β-processes go out of equilibrium (see sect. **2**), so that the ratio of neutron and proton concentrations freezes out.

4) All the local annihilations of antinucleons have finished before the period of nucleosynthesis. Effects of annihilation in the period of nucleosynthesis are negligible.

Parameters:

1) the baryon density ϱ_B,

2) the rate of expansion ξ,

3) the frozen ratio of neutron and proton concentrations n/p,

4) the rate of decay of the neutron C.

The relationship between temperature and density

$$\varrho_b = hT^3 \tag{A.1}$$

was used in the calculations. Here h stands for the entropy. If there were no other additional sources of entropy after the nucleosynthesis, except electron-positron pair annihilation, there would be an unambiguous relationship between the present mean baryon density and the parameter h_0:

$$\varrho_b(T = 2.7 \text{ K}) = 7.15 \cdot 10^{-27} h_0 \text{ g/cm}^3 . \tag{A.2}$$

The rate of isotropic expansion was taken as

$$\frac{1}{V}\frac{dV}{dt} = \xi\sqrt{24\pi G\varrho}\,, \tag{A.3}$$

where V is the relative velocity between the two objects, ϱ is the cosmological density, G is the gravitational constant and ξ is a parameter accounting for the possible presence of new kinds of particles in the Universe. The general limit on ξ from the results of nucleosynthesis was first obtained by SHVARTSMAN [94]. Within the framework of GUTs, the parameter ξ characterizes the number of generations of quarks and leptons, measuring the number of different types of neutrino.

In calculations [78, 79] the lifetime of the neutron relative to the β-decay was taken as

$$\tau_n = 926/C \text{ s} . \tag{A.4}$$

The parameter C accounts for the ambiguities of the experimental value of the neutron lifetime. The neutrino degeneracy [79, 90] in the period of nucleosynthesis may influence the rate of β-decay either.

In the standard model the value $\xi = C = 1$ was accepted. The change of entropy after nucleosynthesis was assumed to be related to the electron-positron pair annihilation only. In the case of additional sources of entropy (cf. [91]) relationship (A.2) between the modern baryon density and the parameter h_0 is to be modified.

Results. The results of calculations are presented in fig. A.1-A.3 taken from [78, 79, 82]. In fig. 9.3 the dependence of the abundances of light elements in the standard ($\xi = C = 1$) model on the baryon density is given. Remind

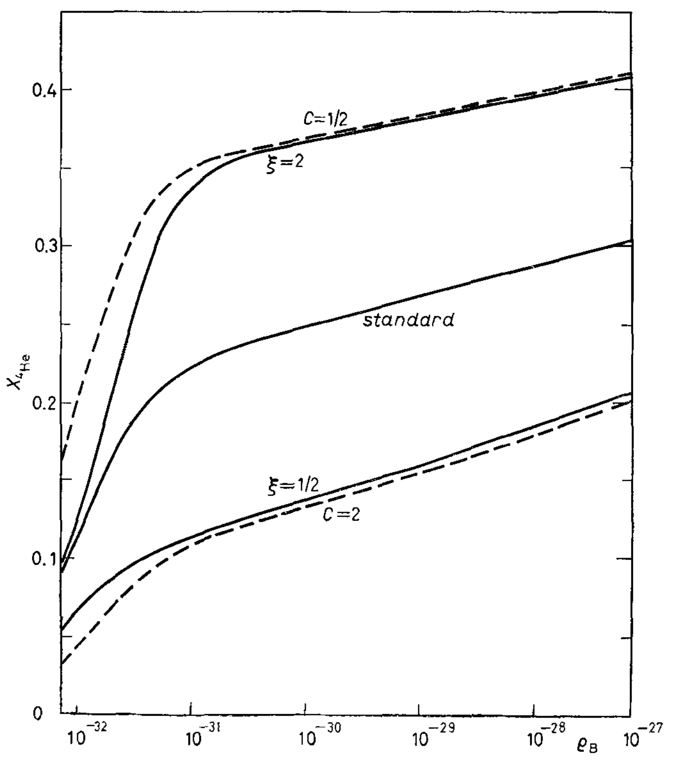

Fig. A.1. – The deviations of the primordial ^{4}He mass fraction for different baryon densities ϱ_b from the predictions of the standard model.

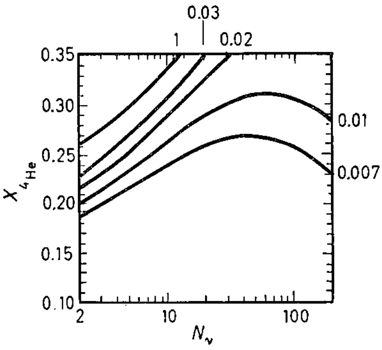

Fig. A.2. – The dependence of the ^{4}He mass fraction for different baryon densities ϱ_b on the number of neutrino species.

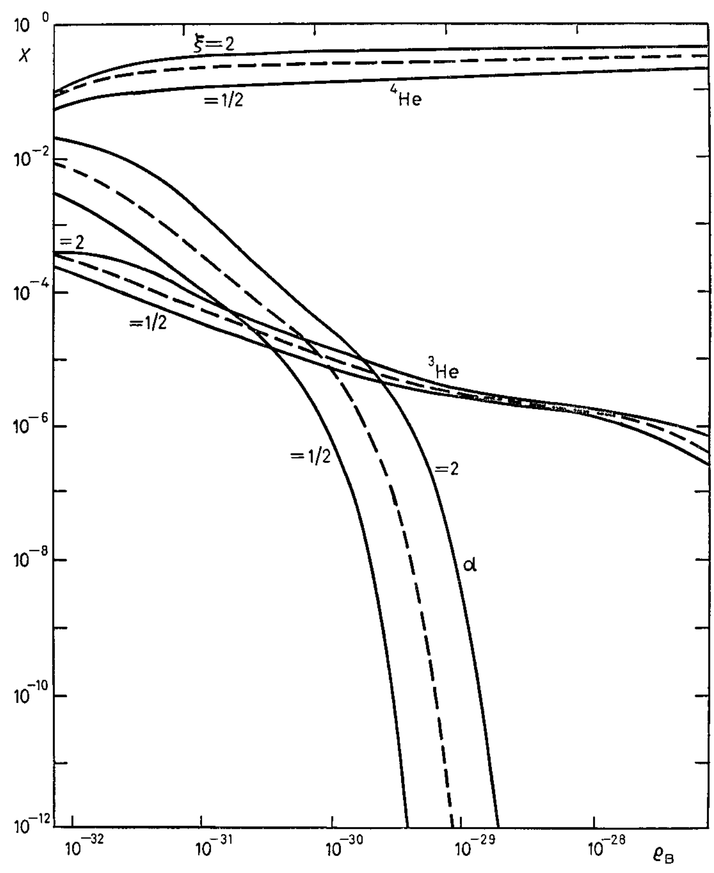

Fig. A.3. – The deviations of d, ^{3}He, ^{4}He mass fractions from the predictions of the standard model (dashed lines) at $\xi = \frac{1}{2}$ and $\xi = 2$.

(see sect. **2**) that the critical density (for $H = 50$ km/s·Mpc) is $\varrho_c = 5 \cdot 10^{-30}$ g/cm^3. The mean density of the visible (luminous) matter (see sect. **2**) is $1.5 \cdot 10^{-31}$ g/cm^3. The output of deuterium depends strongly on the baryon density. We see from fig. 9.3 that at $\varrho_b \sim 10^{-30}$ g/cm^3 a deficit of the predicted deuterium abundance as compared to observations arises. The helium-4 output is dependent on the baryon density at very small ϱ_b only.

The primordial ^{4}He abundance is a rather good indicator of $\xi \neq 1$ and $C \neq 1$. In fig. A.1 the output of ^{4}He is given for $\frac{1}{2} < \xi < 2$ and $\frac{1}{2} < C < 2$. A comparison with the observations of the ^{4}He abundance provides limits on the possible deviations of ξ and C from 1.

The dependence of the ^{4}He abundance on the number of species of relativistic particles for various baryon densities is given in fig. A.2. In fig. A.3 the dependence of d, ^{3}He and ^{4}He abundances on ξ is given.

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