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BARYON ASYMMETRY IN INFLATIONARY UNIVERSE

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A b s t r a c t

It is shown that the baryon asymmetry in inflationary universe under certain constraints on masses of superheavy bosons can be larger than that in the standard baryosinthesis scenario. An important property of the model considered is that the final baryon asymmetry does not depend on initial conditions in the early universe in contrast to what occurs in the standard scenario based on (B-L) conserving GUTs.

A new approach to the long-standing cosmological problems such as the horison and flatness problems and the problem of spatial homogeneity and isotropy has been proposed recently [1-5]. This is done in the framework of the so-called inflationary universe scenario in which the universe is supposed to expand exponentially, like de Sitter space for a sufficiently long period. During this period the stress tensor is dominated by the vacuum (cosmological) term.

Inflationary universe scenario is based on the theory of the phase transitions [6] in grand unified theories (GUTs). If the transition from symmetric to asymmetric state is the first order phase transition with a large supercooling, then the stress tensor in the supercooled phase is indeed dominated by the cosmological term $g_{\mu\nu}\rho^{vac}$ and the universe inflates exponentially fast.

The original scenario of this type proposed by Guth [1] results however in a very inhomogeneous and anisotropic universe [7]. The new inflationary universe scenario suggested in ref. [2] (see also refs. [3-5]), seems to be free from these shortcomings. According to this scenario the inflation proceeds during the early stage of the phase transition in the asymmetric phase when the vacuum condensate of the Higgs field φ is not yet completely developed. If the classical field φ grows to its equilibrium value φ_0 slowly enough, then at the end of the de Sitter stage the universe becomes almost exactly flat, homogeneous, and isotropic independently of the initial conditions at the beginning of the phase transition [2].

This scenario provides also a possible solution of the primordial monopole problem [8] and may explain the origin of inhomogeneittles which are necessary for the galaxy formation [9].

It seems to us that the oddds for this scenario to be true might be rather high, However many problems are to be solved before the final version of this scenario is formulated. One of the most interesting and important problems here is the problem of the baryosynthesis [10] in the new inflationary universe scenario. In the present paper we shall show that the mechanism of the baryon production in this scenario differs considerably from the usual one.

The final baryon asymmetry depends on the relations between masses of superheavy bosons (vector as well as scalar ones) and can be about two orders of magnitude larger than in the standard scenario.

The particle production in the discussed model looks as follows. For some period of time the de Sitter law of expansion is valid. During this period $\varphi \lesssim H \ll \varphi_0$ and the stress tensor is dominated by the cosmological term

$$T_{\mu\nu} \simeq g_{\mu\nu} \rho^{vac} \tag{1}$$

since the usual matter disappears because of long and fast expansion. The scale factor $a(t)$ behaves as

$$a \sim \exp(Ht), \quad H = \left(\frac{8\pi \rho^{vac}}{3 M_P^2}\right)^{1/2} \tag{2}$$

where $M_P = 1.2 \cdot 10^{19}$ GeV is the Planch mass.

After the exponential expansion period (which, hopefully, lasts for $\Delta t \gtrsim 60\ H^{-1}$) the classical field $\varphi(t)$ rapidly grows and then convergently ascillates near some constant equilibrium value φ_0 approximately as

$$\varphi(t) \simeq \varphi_0 (1 - e^{-\gamma t} \cos m_\varphi t) \qquad (3)$$

where m_φ is the mass of the Higgs boson φ when the classical field φ becomes equal to φ_0 .

The oscillation damping factor γ for the noninteracting field φ would be given by 3/2H due to the "friction" term 3H in the equation for the spatially homogeneous field φ in the expanding universe:

$$\ddot{\varphi} + 3H\dot{\varphi} + m^2(\varphi)\varphi = 0 \qquad (4)$$

However in our case at the stage of rapid oscillations (3) (though not at the exponential expansion stage) the main contribution to γ arises because of the particle production by the oscillating field $\varphi(t)$ (3).

To get insight into this process one should consider a concrete theory in which the new inflationary scenario can be realized. We will discuss here the SU(5) Coleman-Weinberg theory with the symmetry breaking SU(5) — SU(3) x SU(2) x U(1) due to appearance of the classical field $\Phi = \sqrt{2/15}$. $\varphi_0 (1, 1, 1, -3/2, -3/2)$. Numerical values of the essential parameters in this theory are the following: vector meson masses, $M_X \simeq 4 m_\varphi$; equilibrium value of the Higgs field, $\varphi_0 \approx 3 M_X$; vacuum energy, ρ^{vac} = $= 9M_X^4 / 32\pi^2 = m_\varphi^2 \varphi_0^2 / 16$; the Hubble constant at the beginning of the phase transition, $H_0 = \sqrt{3/\pi}\, M_X^2 / (2M_\mathfrak{p})$. The typical value of M_X is $5 \cdot 10^{14}$ GeV, and so H_0 =

$\approx 10^{10}$ GeV.

The theory of particle production by the large oscillating field is rather complicated and will not be considered here. Fortunately, however, one may obtain reasonable estimates of the particle production rate by means of perturbation theory. To that end one should calculate the lowest-order contribution of particles of any given type to the imaginary part of the effective action $V(\varphi)$ in external field (3). This gives in particular the lowest-order estimate of the probability of the φ particle production per unite time per unit volume

$$\frac{dW_\varphi}{dV\,dt} = a\,\frac{m_\varphi^4}{512\pi} \tag{5}$$

where a is some numerical factor, $a = O(1)$.
An important feature of this result is the m^4 – dependence of $dW/dV dt$ on the mass of the particles produced. Therefore if the φ -bosons are the heaviest Higgs bosons in the theory (as we shall admit for a while), then only these Higgs bosons will be produced by the field $\varphi(t)$. As for the vector bosons X, Y, they are not produced in the lowest order in the gauge coupling g^2 since M_X is (approximately) four times larger than the frequency m_φ of the oscillations (3). However these particles are produced in higher orders in g^2, and in fact this leads to the contribution to $dW/dVdt$ of the same order in g^2 since the external field is large, $\varphi \approx \varphi_0 \approx M_X/g$. Nevertheless some suppression of X, Y – production appears due to

combinatorial factors. On general grounds one would expect $dW_{X,Y}/dVdt \sim M_X^4 e^{-\ell M_X/m_\varphi}$, where $\ell = O(1)$, and indeed the simplest diagrams for $V(\varphi)$ which give a nonvanishing contribution to $dW_{X,Y}/dVdt$ are $O(g^{20})\varphi^{20}M_X^{-16}$, and the corresponding contribution to $dW_{X,Y}/dVdt$ is

$$\frac{dW_{X,Y}}{dV\,dt} = C\cdot 1.5\cdot 10^{-4}\cdot 36\,\frac{M_X^4}{512\pi} \approx C\cdot 6\cdot 10^{-3}\,\frac{M_X^4}{512\pi} \qquad (6)$$

where $C = O(1)$, $1.5\cdot 10^{-4}$ is the above-mentioned combinatorial factor and 36 is the number of heavy vector meson species (including helicity states) in the SU(5) theory.

Taken by their face values expressions (5) and (6) predict a somewhat larger production of heavy vector mesons with $M_X = 4m_\varphi$ as compared to the production of scalar particles with $m \leq m_\varphi$. There can be however some enhancement of scalar particle production if $m = m_\varphi$ (i.e. quanta of φ-field itself) because of resonant effects. At the moment we cannot definitely say what kind of particles, the ones with $m = m_\varphi$ or with $m > m_\varphi$, are dominately produced by the oscillating field $\varphi(t)$. This problem deserves a further investigation and in what follows we consider both possibilities: 1) the field $\varphi(t)$ transforms predominantely into X, Y-bosons ("vector dominance"); 2) the classical field $\varphi(t)$ transforms predominantely into φ bosons ("scalar dominance").

1) Let us consider first the "vector dominance" scenario and estimate the time during which the vacuum energy $\rho^{vac} = 9M_X^4/32\pi^2$ transforms into the energy of non-relativistic X, Y bosons:

$$\Delta t \simeq \frac{\rho^{vac}}{M_X \frac{dW}{dVdt}} \simeq 5\cdot 10^3 \frac{M_X}{M_\varphi}\cdot H_0^{-1} \qquad (7)$$

Note, that for $M_X \lesssim 5\cdot 10^{14}$ GeV the rate of the heavy meson production is higher than the rate of the universe expansion, $\Delta t \ll H_0^{-1}$. In the theories with large M_X creation of particles by the oscillating field $\varphi(t)$ proceeds during the time $\Delta t \gg H_0^{-1}$, which would modify to some extent the baryosynthesis scenario to be discussed. However in the present paper we restrict ourselves for simplicity to the discussion of the theories with sufficiently small M_X, so that $\Delta t \lesssim H_0^{-1}$ (7).

The produced X and Y -bosons decay into quarks and leptons with the lifetime which is also smaller than the characteristic expansion time H_0^{-1}. The baryon asymmetry generated by these decays is

$$\beta_0 = \frac{N_q - N_{\bar{q}}}{N_q + N_{\bar{q}} + N_\ell + N_{\bar{\ell}}} \approx \frac{3}{4}\,\frac{\Gamma(X\to qq) - \Gamma(\bar{X}\to \bar{q}\bar{q})}{\Gamma(X\to all)} \qquad (8)$$

This value is about two orders of magnitude larger than that obtained in the standard scenario (if C and CP violating effects in vector and scalar boson decays are of the same order of magnitude, and in contrast to the standard scenario is born by the nonequilibrium decays of vector but not scalar (Higgs) bosons.

Baryon asymmetry β_0 (8) can be somewhat washed out in the course of the further evolution of the universe. The dangerous processes are the baryonic charge nonconserving inverse decays of the vector X, Y -bosons and the scalar fractionally charged Higgs particles $H^{\pm 1/3}$. To evaluate their destructive role we take into account that the number density of X and Y bosons just after their production is

$$N_X = \frac{\rho^{vac}}{M_X} = \frac{9M_X^3}{32\pi^2} \qquad (9)$$

and the number density of fermions born by their decays is $N_F = 2N_X$. The important point is that this number density is the equilibrium number density corresponding to ρ^{vac} if the number of particle species is about 10^2 (as is indeed the case). Consequently the thermodynamic equilibrium is established by the processes $2 \leftrightarrow 2$ which do not change the total particle number and are fast as compared to H_c^{-1}. The resulting plasma temperature is $T_0 \approx m_X/5$, and thus the inverse vector mesons decays cannot considerably diminish baryon excess (8). The processes with $H^{\pm 1/3}$ seem to be more important. Their contribution to washing out the initial baryon asymmetry can be evaluated as $\exp\{-(\Gamma_H/H_0)\cdot(T_0/m_H)^2\cdot const\}$, where $const = O(1)$, $\Gamma_H = \alpha_H m_H/2$ is the decay width of $H^{\pm 1/3}$ and $\alpha_H \approx 10^{-4}$. This does not considerably change our result (8) if $H^{\pm 1/3}$ are heavy enough ($m_H \gtrsim 5\cdot 10^{13}$ GeV). Otherwise the initial value of β_0 is completely washed out and we come to the standard scenario of baryon asymmetry generation through the near-equilibrium decays of $H^{\pm 1/3}$-mesons.

2) If $\varphi(t)$ produces mainly φ-bosons the baryon generation becomes more complicated. The number density of φ bosons just after the phase transition is

$$N_{\varphi_0} = \frac{m_\varphi \varphi_0^2}{16} \tag{10}$$

Then the universe becomes Friedman one with the energy momentum tensor dominated by nonrelativistic φ bosons:

$$\rho(t) = \frac{M_\varphi^2}{6\pi (t+t_0)^2} \tag{11}$$

where t_0 is defined by $\rho(0) = M_\varphi^2 / 6\pi t_0^2 = m_\varphi \varphi_0^2 / 16$.

Note that the number density of φ bosons is much larger than the equilibrium number density which corresponds to the energy density $\rho(0)$.

The further development of the initial φ -boson state depends upon the relation between their decay rate into a pair of lighter Higgs bosons χ

$$\Gamma_\varphi = \Gamma(\varphi \to \chi\bar{\chi}) = \frac{m_\chi^4}{4\pi m_\varphi \varphi_0^2} \left(1 - \frac{4m_\chi^2}{m_\varphi^2}\right)^{1/2} \tag{12}$$

and the rate of their number density decrease through the processes $3\varphi \to 2\varphi$, $4\varphi \to 2\varphi$ and so on

$$\Gamma(3\varphi \to 2\varphi) = \frac{N_\varphi^2}{(2m_\varphi)^3} |A(3\varphi \to 2\varphi)|^2 \approx \frac{N_\varphi^2}{(2m_\varphi)^3} \left|\frac{\lambda^2 \varphi_0}{m_\varphi^4}\right|^2 \approx 0.2 \frac{m_\varphi^3}{\varphi_0^2} \tag{13}$$

The rates of the processes $4\varphi \to 2\varphi$ etc are much smaller.

Comparing expressions (12) and (13) one can conclude that at the first stage of the process the quasitermalization of the φ -boson gas takes place which leads to their equilibrium number density and the equilibrium distribution in energy, particles of other types being not produced and only later the φ -decays enter the game. This picture is valid if the rate of the universe expansion (see eqs. (2), (4), and (11)) is smaller than $\Gamma(3\varphi \to 2\varphi)$. Their ratio is given by.

$$\frac{H_0}{\Gamma(3\varphi \to 2\varphi)} \simeq \frac{3\varphi_0^3}{m_\varphi^2 M_p} \qquad (14)$$

As it was already noted, $\varphi_0 \simeq 3M_X \simeq 12 m_\varphi$, so the expansion is slow if $m_\varphi < 5\cdot 10^{-3} M_p$, which is true in the standard SU(5)-model. Thus the energy density of φ's remains practivally unchanged, $\rho = m_\varphi^2 \varphi_0^2/16$, while their number density decreases ans becomes equal to

$$N_{\varphi 1} = 0{,}12 \left(\frac{30}{16\pi^2}\right)^{3/4} (m_\varphi \varphi_0)^{3/2} \simeq 0.2 N_{\varphi 0} \qquad (15)$$

The corresponding temperature is $T \simeq 2m_\varphi$.

Note that despite the cannibalism of φ -bosons their number density remains greater than the equilibrium one because particles of other types have not enough time to be produced at this stage of the process. This results in a larger baryon asymmetry because the number density of quarks which are produced after the successive decays $\varphi \to H^{+1/3} + H^{-1/3}$ and $H^{\pm 1/3} \to 2q$, is also larger than the equilibrium one.

If the fractionally charged $H^{\pm 1/3}$ bosons are the nearest in mass to φ then the decay $\varphi \to H^{1/3} + H^{-1/3}$ is the dominant one. In other case, when there exist scalar mesons χ with a mass such that $m_\varphi > 2m_\chi > 2m_H$, the process of quark production becomes more complicated:

$\varphi \to \chi\bar{\chi} \to 2(H\bar{H}) \to 4(q\bar{q})$.

Let us first consider the more simple case when $m_H > m_\chi$. If other processes can be neglected, then after the chain of decays $\varphi \to H + \bar{H}$, $H \to qq, \bar{q}\bar{\ell}$, $\bar{H} \to \bar{q}\bar{q}, q\ell$ the quark-lepton plasma is produced with the relative baryon excess:

$$\beta_0 = \frac{N_q - N_{\bar{q}}}{N_q + N_{\bar{q}} + N_\ell + N_{\bar{\ell}}} = \frac{3}{4}\,\frac{\Gamma(H \to qq) - \Gamma(\bar{H} \to \bar{q}\bar{q})}{\Gamma(H \to all)} \qquad (16)$$

This number is larger than that obtained in the standard scenario [11]

$$\tilde{\beta} = \frac{1}{K_{tot}}\,\frac{m_H^2 K_{tot}^{1/2}}{\Gamma_H m_\varphi}\,\beta_0 \approx \left(10^{-2} \div 10^{-3}\right)\beta_0 \qquad (17)$$

where $K_{tot} \simeq 10^2$ is the total number of particle species (including the number of spin states) in the equilibrium plasma.

However the comparison of the decay rate $\Gamma_\varphi = \Gamma(\varphi \to H\bar{H})$ (12) with the expansion rate $H_0 = \sqrt{\pi/6}\, m_\varphi \varphi_0 M_{Pl}^{-1}$, that is

$$\frac{\Gamma_\varphi}{H_0} \simeq 0.1\,\frac{m_H^4 M_{Pl}\sqrt{1 - 4m_H^2/m_\varphi^2}}{m_\varphi^2 \varphi_0^3} \approx 10^{-4}\,\frac{m_H^4 M_{Pl}}{m_\varphi^5}\sqrt{1 - \frac{4m_H^2}{m_\varphi^2}} \qquad (18)$$

shows that $\Gamma_\varphi \ll H_0$ if m_H is sufficiently small. In this case the decay of nonrelativistic φ-bosons leads to the entropy increase which is proportional to $(H_0/\Gamma_\varphi)^{1/2}$ and to the corresponding decrease of β. Finally the following expression holds for β after the total thermalisation:

$$\beta \approx \beta_0 \left(\frac{\Gamma_\varphi}{H_0}\right)^{1/2} \cdot 10 K_{tot}^{-1/4} \left(\frac{N_{\chi_1}}{m_\varphi^3}\right)^{1/4} \qquad (19)$$

where β_0 is given by eq. (16) and N_{χ_1} is given by eq. (15). The factor $10 K_{tot}^{-1/4} (N_{\chi_1}/m_\varphi^3)^{1/4}$ which leads to some increase of β appears due to the fact that the initial number density of quarks is greater than the equilibrium one. To make an estimate of the final baryon asymmetry in this scenario let us take $m_H \simeq 0.2\, m_\varphi \approx 2\cdot 10^{13}$ GeV. In this case eq. (19) yields $\beta \simeq 0.5\beta_0 \gg \tilde{\beta}$, i.e. the resulting baryon asymmetry is much greater than that in the standard scenario.

In deviation of eq. (19) we have implicitely assumed that H bosons disappear only because of the decays of the type $H \to q + q$, but not because of the annihilation $H + \bar{H} \to 2V$, where V are light vector mesons. Let us analyse, whether this assumption is true. The decay rate is equal to

$$\Gamma_H = \Gamma(H \to 2q) = \frac{\alpha_H m_H}{2} \qquad (20)$$

where $\alpha_H \approx 10^{-4}$ is the coupling constant of coloured Higgs mesons with fermions. With a reasonable choice of the parameters ($m_H \simeq 0.2 m_\varphi$) it appears that $\Gamma_H > \Gamma_\varphi$ In this case H-mesons produced in the decays of φ are relativistic at the moment of the decay $H \to 2q$, and consequently the H-decay rate is $(m_H/m_\varphi)\Gamma_H$.

Taking into account only the processes $\varphi \to H + \bar{H}$ and $H \to 2q, \bar{q}\bar{\ell}$ it is easy to calculate the concentration of these particles as functions of time, the condition $m_H\Gamma_H/m_\varphi > \Gamma_\varphi$ being assumed.

$$
\begin{aligned}
&\tau_\varphi = \exp(-\Gamma_\varphi t), \quad \tau_H = (m_H\Gamma_H)/(m_\varphi\Gamma_\varphi)\cdot\exp(-\Gamma_\varphi t) \\
&\tau_q = \frac{3}{2}\left[1 - \exp(-\Gamma_\varphi t)\right]
\end{aligned}
\tag{21}
$$

where $\tau_j = N_j a^3(t)$ and $a(t)$ is the scale factor with the normalization condition $a^3(0)N_{\varphi_1} = 1$.

Now we can evaluate the concentration of vector bosons produced in the process $H + \bar{H} \to 2V$ as

$$
\begin{aligned}
\tau_V &= v\sigma(H\bar{H} \to 2V)\cdot 2\ln 2\, N_{\varphi_1} t_0^2 \Gamma_\varphi^2 \Gamma_H^{-1} = \\
&= 10^{-4}\frac{\alpha^2}{\alpha_\varphi}\frac{M_\varphi^2 m_H^5}{m_\varphi^3 \varphi_0^4}\left(1 - \frac{4m_H^2}{m_\varphi^2}\right)
\end{aligned}
\tag{22}
$$

which is valid for $\Gamma_H t_0 = \sqrt{6/\pi}\,\alpha_H m_H M_\varphi m_\varphi^{-1}\varphi_0^{-1}$ (here $\alpha = 2\cdot10^{-2}$). One can see that the vector meson production is small for sufficiently small m_H ($m_H \simeq 0,\ m_\varphi$).

Analogously the contribution of other processes of particle production or the inverse $H^{\pm 1/3}$-decays does not considerably change our result (19) because at the essential stage the expansion rate is large in comparison with Γ_φ.

In a more complicated case when $m_\varphi > 2m_\chi > 2m_H$

a somewhat smaller value of β can be obtained. The concrete number depends on the ratio of expansion rate and different decay rates and can be evaluated along the same lines as above. The baryon asymmetry could be very much suppressed in a rahter improbable case when $m_H \gg m_y$.

We see therefore that the baryon asymmetry generation in the inflationary universe crucially depends on the parameters of the theory, and for certain relations between masses of superheavy bosons the resulting baryon asymmetry proves to be one or two orders of magnitude greater than that in the standard baryosynthesis scenario.

There is also a rather important point concerning the dependence of the baryon asymmetry on initial conditions. It is believed usually [10] that in the standard scenario any initial baryon asymmetry is washed out by the B -non-conserving processes during thermodinamically equilibrium stage. However this might be true only in the theories in which strong B-L violation is possible, and it is definitely not the case for the simple models considered in the literature, in which B-L is conserved. Only in the inflationary universe scenario, when practically no matter is left after the exponential expansion, all the initial conditions in the early universe become unimportant for its evolution after the phase transition, and the final baryon asymmetry becomes a calculable number. This independence of the present baryon asymmetry on the initial conditions in the universe is a very attractive feature of the inflationary universe scenario.

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