Late Reheating, Hadronic Jets, and Baryogenesis

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If inflaton couples very weakly to ordinary matter, the reheating temperature of the Universe can be lower than the electroweak scale. In this Letter we show that the late reheating occurs in a highly nonuniform way, within narrow areas along the jets produced by ordinary particles originated from inflaton decays. Depending on inflaton mass and decay constant, the initial temperature inside the lumps of the overheated plasma may be large enough to trigger the unsuppressed sphaleron processes with baryon number nonconservation. This allows for efficient local electroweak baryogenesis at reheating temperatures $T_R \sim \mathcal{O}(10)$ GeV.

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Introduction.—The inflationary paradigm (for a review see, e.g., books of Ref. [1]) happened to be very successful for understanding the basic properties of the universe. It is assumed that the energy density of the early universe was dominated by a potential energy of a scalar field—inflaton. The accelerating expansion of the universe then leads to a solution of the horizon and flatness problems, whereas quantum fluctuations of the inflaton field result in density perturbations necessary for structure formation and seen as the temperature fluctuations of the cosmic microwave background radiation.

The exponential expansion of the universe during inflation must be eventually replaced by a radiation dominated epoch, which started at least somewhat before nucleosynthesis and lasted till about recombination. The process of transfer of the energy density of the inflaton to ordinary matter is usually called reheating, and several scenarios for how it may proceed have been proposed. Quantitatively, they can be distinguished by the value of the reheating temperature T_R, below which the universe expansion is dominated by radiation. If the coupling of inflaton to the ordinary matter is sufficiently strong, the energy transfer occurs very rapidly due to the phenomenon of broad parametric resonance right after the inflationary stage [2]; this leads to high reheating temperatures $T_R \sim 10^{10}$ GeV or so. If, on the contrary, the coupling is very weak, the exponential expansion of the universe is first replaced by a matter dominated period during which inflaton oscillates without dissipation. Then, the perturbative decays of inflaton heat the universe up to some temperature T_R. In the latter case the reheating temperature can be very low, with the only reliable bound $T_R > 1$ MeV coming from the successful predictions of the big bang nucleosynthesis. The small values of the reheating temperatures can naturally occur in certain supergravity models (see, e.g., [3]) where the decay rate of the inflaton is suppressed by the Planck scale.

A successful cosmological model should also explain the absence of antimatter in the universe and the baryon to entropy ratio $n_B/s \simeq 9 \times 10^{-11}$ [4]. The baryogenesis must occur after inflation since otherwise all created baryonic excess will be exponentially diluted.

Quite a number of different baryogenesis mechanisms exist in theories with high reheating temperatures. The explosive particle production during inflaton decay in the wide resonance case [5] may lead to production of grand unified theory leptoquarks, the subsequent CP-violating and baryon number nonconserving decay of which gives rise to baryon asymmetry of the universe. The gravitational production of superheavy particles has a similar effect [6]. The thermal [7] or nonthermal [8] production of heavy Majorana neutrinos may prepare suitable initial conditions for generating the lepton asymmetry [9], which is transformed then to baryon asymmetry due to anomalous electroweak number nonconservation [10]. If the reheating temperature is greater than the electroweak scale, an electroweak baryogenesis (for reviews see [11]) can take place.

The theories with high T_R may, however, be in conflict with observations because of overproduction of dangerous relics like gravitinos [12]. From this point of view the theories with small reheating temperature are more advantageous, as the production of unwanted particles is automatically suppressed. At the same time, the problem of baryon asymmetry of the universe is much more difficult if T_R is relatively small, simply because in this case the low energy baryon number nonconservation is required. Thus, none of the mechanisms related to the production of heavy leptoquarks or Majorana neutrinos is operative, and one has to rely on some variant of electroweak [13,14] or Affleck-Dine [15] baryogenesis.

The electroweak baryogenesis, occurring in expanding and almost equilibrium plasma, is highly constrained, as it requires the freezing of the sphaleron processes after the first order phase transition. This condition can be

converted in the upper bound on the Higgs mass in the minimal standard model [16]. This bound cannot be satisfied with the experimental value of the top quark mass [17], so that new physics is required. In the minimal supersymmetric standard model there exists a (small) region of the parameter-space leading to a sufficiently strong first order phase transition (see [18] and references therein).

In [14] it was pointed out that the maximum temperature during reheating can, in fact, be much larger than the temperature T_R and that the rate of the universe expansion in the region of the electroweak phase transition can be faster than it is usually assumed. This allowed the authors to relax somewhat the Higgs mass bound for the electroweak baryogenesis.

In this Letter we show that the late inflaton decays heat up the plasma in a very nonuniform way and that the *local* temperature along the trajectories of decay products is substantially higher than the average one. For a range of parameters the mean temperature of the plasma is *small* enough [$T \sim \mathcal{O}(10)$ GeV] to shut off the sphaleron transitions, whereas the temperature of the overheated regions is *large* enough [$T \sim \mathcal{O}(100)$ GeV] to switch them on. Electroweak baryogenesis is then possible in these regions. The overheated regions cool down due to diffusion and expansion. This process has a highly nonequilibrium character, so no first order phase transition is required to satisfy the washout condition and thus no bound on the Higgs mass is implied at all. Remarkably, the resulting baryon asymmetry does not depend much on many details of the process and may be consistent with the observed one for large enough CP violation.

Local overheating.—Let M_ϕ be the inflaton mass and $\Gamma_\phi = f_\phi M_\phi$ its width. Assuming the instantaneous decay of the inflaton at time $t_\phi = 1/\Gamma_\phi$ the reheating temperature is given by $T_R = \sqrt{f_\phi M_\phi M_0}$, where $M_0 = M_{\rm Pl}/1.66 g_*^{1/2}$, $M_{\rm Pl}$ is the Planck mass, $g_* \sim 100$ is the effective number of massless degrees of freedom. The typical numbers we are interested in are $T_R < T_W \sim 100$ GeV (here T_W is the freezing temperature of the sphaleron processes), and $M_\phi > 10^{10}$ GeV, which requires a rather weak decay constant $f_\phi \sim 10^{-24}$, or smaller. We shall assume, for simplicity, that inflaton decays into a quark-antiquark pair, though other decay channels lead essentially to the same result.

The number density of inflatons decreases with time as $n_\phi a^3 = n_0 \exp(-\Gamma_\phi t)$ with scale factor a. The first decays occur essentially in the vacuum, whereas at $t \sim t_\phi$ inflatons are surrounded by the plasma with the temperature $T \sim T_R$. The decay products of inflaton, ultrarelativistic quarks with the energy $M_\phi/2 \gg T_R$, are injected into the plasma and heat it locally. Our aim now is to understand the typical size and geometry of the overheated regions, as well as their temperature.

The dynamics of energy losses of high energy quarks and gluons in quark-gluon plasma is a complicated problem which must incorporate the Landau-Pomeranchuk-Migdal effect [19] and non-Abelian character of interaction of quarks and gluons. The main features of it have been understood only recently (see [20] and references therein), the crucial one being that the energy loss per unit length is increasing in infinite plasma as a square root of parton energy; see Eq. (2).

The most significant part of energy losses is related to the soft gluon emission in the multiple scattering of the hard parton on the particles of the plasma. The energy spectrum I of the emitted gluons per unit length has the following approximate form for $\omega_{\rm BH} \ll \omega \leq E_0$:

$$\omega \frac{d^2 I}{d\omega dz} = \frac{2}{\pi} \frac{\alpha_s}{\lambda_g} \sqrt{\frac{\omega_{\rm BH}}{\omega} \ln\left(\frac{\omega}{\omega_{\rm BH}}\right)}, \tag{1}$$

where ω and E_0 denote the gluon and parton energies, α_s is the QCD coupling constant, λ_g is the mean-free path of the gluon, and $\omega_{\rm BH}$ is the Bethe-Heitler (BH) frequency, $\omega_{\rm BH} = \lambda_g \mu^2$. Here μ is the typical screening mass which is assumed to be of the order of the Debye mass in the plasma. From (1) one gets the stopping distance of the initial parton $L_{\rm tot}$:

$$L_{\rm tot} = \frac{\pi}{2} \frac{\lambda_g}{\alpha_s} \sqrt{\frac{E_0}{\tilde{\omega}_{\rm BH}}}, \tag{2}$$

where $\tilde{\omega}_{\rm BH} \simeq \omega_{\rm BH} \log(E_0/\omega_{\rm BH})$ for $E_0 \gg \omega_{\rm BH}$. The average energy of the emitted gluons is $\langle\omega\rangle = \sqrt{\omega_{\rm BH} E_0}/2$, and their number is given by $N_g = 2\sqrt{E_0/\omega_{\rm BH}}$.

If the emitted gluons have energies $\omega \gg \omega_{\rm BH}$, they lose their energies as the parent parton does. This cascade process terminates when the energy of the emitted gluons becomes comparable to or smaller than $\omega_{\rm BH}$. We find the energy for the nth gluons as

$$\langle\omega\rangle_n = \frac{1}{4^{1-1/2^n}} \omega_{\rm BH} \left(\frac{E_0}{\omega_{\rm BH}}\right)^{1/2^n}, \tag{3}$$

so number $N_{\rm BH}$ of the cascade steps from $\langle\omega\rangle_0 = E_0$ to $\langle\omega\rangle_{N_{\rm BH}} \leq \omega_{\rm BH}$ is typically 3–4 for $E_0 \sim 10^{10}$–10^{11} GeV.

To find the geometry and volume of the region where the emitted gluons deposit their energy we note that the radiation with the frequency ω is mainly concentrated in the cone with a small angle θ, which can be computed with the help of the results of [21] and is given by

$$\theta_\omega^2 \simeq \frac{L_\omega \overline{\omega}_{\rm BH}}{\lambda_g^2 \omega^2}, \tag{4}$$

where L_ω is the length of the energy loss for the gluon with energy ω,

$$L_\omega = \frac{2\pi}{9} \frac{\lambda_g}{\alpha_s} \sqrt{\frac{\omega}{\overline{\omega}_{\rm BH}}}, \tag{5}$$

$\overline{\omega}_{\rm BH} = \omega_{\rm BH} \ln(\omega/\omega_{\rm BH})$. The transverse distance $r_\omega = L_\omega \theta_\omega$ traveled by a gluon with the energy ω is then

$$r_\omega \simeq \left(\frac{2\pi}{9\alpha_s}\right)^{3/2} \frac{\lambda_g^{1/2}}{\omega_{\rm BH}^{1/4}\omega^{1/4}} \times \left[\ln\left(\frac{\omega}{\omega_{\rm BH}}\right)\right]^{-1/4} \quad (6)$$

(the log factor should be neglected at $\omega \sim \omega_{\rm BH}$). It is seen that, although the cascade process produces gluons having various energies, the largest r_ω is due to the gluons with the lowest energy $\omega \simeq \omega_{\rm BH}$. Thus, the overheated region has an approximate cylindrical form with the length given by Eq. (2) and with the radius $r_* \simeq r_{\omega_{\rm BH}}$. Its volume is $V = \pi r_*^2 L_{\rm tot}$.

Now we are at the point to estimate an effective temperature inside the cylinder in which the emitted gluons are living. By assuming the energy conservation and the rapid thermalization, we obtain that the effective temperature T_*, found from $\frac{\pi^2 g_*}{30} T_*^4 = \frac{E_0}{V}$, is given by

$$T_* \simeq 5.3 \times 10^{-2} \left(\frac{100}{g_*}\right)^{1/4} \mu^{3/4} \left[\frac{E_0}{\lambda_g} \ln\left(\frac{E_0}{\omega_{\rm BH}}\right)\right]^{1/8}. \quad (7)$$

For numerical estimates we take the nonperturbative value of the Debye screening mass found in [22]. As for the gluon mean-free path, we take the value found in [23] and denoted as $\lambda_g^{\rm el}$ or simply take a gluon damping rate γ_g from [24]. They are different by roughly one order of magnitude, which allows one to get an estimate of the uncertainties in (7).

For example, when $T_R = M_Z$ and $E_0 = 10^{11}$ GeV, we find (cf. Fig. 1) that the effective temperature is $T_* \simeq 212$ GeV $(\lambda_g = 1/\gamma_g)$ or $T_* \simeq 151$ GeV $(\lambda_g = \lambda_g^{\rm el})$. Figure 1 clearly demonstrates that the local temperature of the overheated regions can substantially exceed the freezing temperature of the sphaleron processes, provided the background cosmic temperature is relatively high, say, $T_R \gtrsim 10$ GeV, and the energy of the parent parton is extremely high, say, $E_0 \gtrsim 10^{10}$ GeV.

Baryogenesis.—Let us assume now that the parameters of the inflaton are such that the temperature T_* is high enough: $T_* > T_{\rm sph} \sim 100\,{\rm GeV} > T_W$. Here $T_{\rm sph}$ is the temperature above which the rate of the sphaleron transitions is unsuppressed and is given by (per unit time and unit volume) $\Gamma_{\rm sph} \simeq \kappa \alpha_W^5 T_*^4$, where $\kappa \sim 10$ [25]. At the same time, the reheating temperature T_R can be small enough, so that the baryon number is conserved away from the overheated regions (i.e., $T_R < T_W$). This highly nonequilibrium situation is possible at any choice of parameters of the underlying electroweak theory, and, therefore, the washout bound of [16] is not applicable.

The baryon asymmetry of the universe can be estimated as the number of sphaleron transitions which take place inside the overheated regions and go asymmetrically due to CP violation:

$$\frac{n_B}{s} \simeq \frac{n_{\rm parton}}{s} \times \kappa \alpha_W^5 T_*^4 V \Delta t \times \delta_{\rm CP}, \quad (8)$$

where $n_{\rm parton}$ is the number density of the highly energetic partons coming from the inflaton decay, Δt denotes the lasting time of the rapid sphaleron processes, and we have introduced $\delta_{\rm CP}$ to represent the effective magnitude of CP violation.

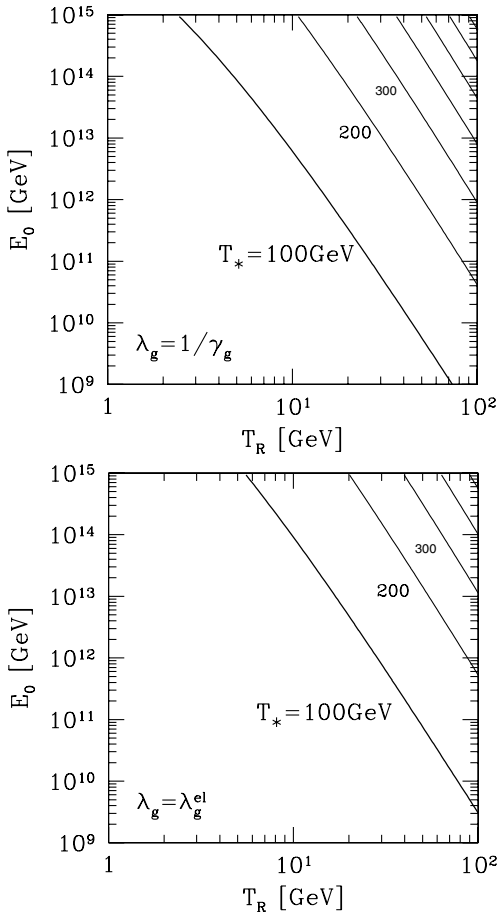

FIG. 1. Contour plots of effective temperature T_* in the T_R-E_0 plane for $\lambda_g = 1/\gamma_g$ (top) and $\lambda_g = \lambda_g^{\rm el}$ (bottom).

With the two-particle decays of the inflaton the number of partons is simply $n_{\rm parton} = 2n_\phi$, and $\frac{n_\phi}{s} \simeq \frac{3}{4}\frac{T_R}{M_\phi}$ from the condition defining the reheating temperature. Putting all factors together, we get

$$\frac{n_B}{s} \simeq 10^{-8} T_R \Delta t \delta_{\rm CP}. \quad (9)$$

Quite amazingly, besides the expected CP-violating factor, the result depends only on the reheating temperature and on the sphaleron transition time. In particular, the temperature of the overheated regions has canceled out from Eq. (9). What is important is that the overheated regions must be in the symmetric phase of the electroweak theory, where baryon number nonconservation is not suppressed.

The overheated regions cool down with the growth of the volume due to diffusion, and the rapid spharelon processes terminate eventually when the temperature inside becomes $T_{\rm sph}$. From the energy conservation the radius of the overheated region at $T_{\rm sph}$ is given by $\ell = r_*(T_*/T_{\rm sph})^2$

and is reached after the diffusion time $\Delta t \simeq \ell^2/4D$ (we find from [26] that the diffusion coefficient $D \sim 1/\gamma_g \sim 0.1\lambda_g^{-1}$), which gives $\Delta t \simeq \frac{r_*^2}{4D}(\frac{T_*}{T_{\rm sph}})^4$. For a region of parameters Δt is long enough to thermalize W bosons which are essential to the sphaleron processes.

Finally, one gets

$$\frac{n_B}{s} \gtrsim (1\text{–}10) \times 10^{-7}\delta_{\rm CP}, \qquad (10)$$

depending on the estimate of the gluon mean-free path discussed above. So, the extension of the standard model with suitable CP violation may work.

Conclusions.—We have shown that a successful electroweak baryogenesis can take place in inflation models with low reheating temperatures $T_W > T_R \gtrsim 10$ GeV. The scenario suggested is of highly nonequilibrium nature, which allows us to overcome the stringent limitations related to the Higgs mass. Though our estimates are rather rude (even parton losses in plasma have rather large uncertainties due to the lack of exact knowledge of the kinetic coefficients), it is clear that the increase of the inflaton mass makes the temperature of the overheated regions higher, which triggers the mechanism at some critical inflaton mass.

Generally speaking, the plasma overheating by the processes discussed above will take place along the trajectories of decay products of any sufficiently heavy particles provided their decay products interact with the background plasma.

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