Asymmetric baryon capture by primordial black holes and baryon asymmetry of the universe

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We have refined our previously suggested scenario of generation of the cosmological baryon asymmetry through an asymmetric capture of baryons and antibaryons by primordial block hole [30]. It is found that in the limit of weak interactions of hypothetical heavy baryons with the primeval plasma the effect can be strongly enhanced and the observed magnitude of the asymmetry can be obtained for a wide range of the model parameters.

I. INTRODUCTION

The generally accepted mechanism of generation of the cosmological baryon asymmetry was suggested by Sakharov [1] in 1967. He formulated the following three necessary conditions for the baryogenesis:

1. Violation of C and CP symmetries in particle physics.
2. Non-conservation of baryonic number $B$.
3. Deviation from thermal equilibrium in the early universe.

Out of these three Sakharov’s principles the nonconservation of baryons remains yet unconfirmed by experiment, so scenarios of baryogenesis which can operate without assumption of baryon nonconservation are certainly of interest.

The magnitude of the asymmetry expressed in terms of the present day number densities of baryons, antibaryons, and photons of the cosmic microwave background radiation is equal to (see e.g. review [2]):

$$\beta = \frac{n_B - n_{\bar{B}}}{n_{\gamma}} \approx 6 \cdot 10^{-10}$$

where $n_B$ and $n_{\bar{B}}$ are respectively the number densities of baryons and antibaryons (note that today $n_{\bar{B}} \ll n_B$), $n_{\gamma} = 411(T_{\gamma}/2.73^oK)^3$ cm$^{-3}$, and $T_{\gamma} = 2.73^oK$ is the present day temperature of the cosmic microwave background (CMB) radiation.

There are many different models of baryogenesis which with properly chosen parameters lead to the correct value of the asymmetry, for reviews see refs. [3–9].

In our scenario excessive antibaryons are hidden inside black holes, so one may say that baryonic number is formally conserved. However, low mass black holes quickly evaporate and disappear from our world without any trace. Thus baryon number is not conserved and the global $U(1)$ symmetry associated with baryonic number is broken. The breaking of all global symmetries, including that related to baryon number, by black holes is known for a long time. In particular, this breaking can lead to proton instability [10]. Though the life-time with respect to this decay is fantastically long, $\tau_p \sim 10^{45}$ years, one has to admit that $B$-conservation in particle physics does not save the proton life. In the TeV gravity [11] proton might decay almost instantly but spin and electric charge of proton could suppress the virtual BH formation and very strongly increase the proton life-time [12].

The idea that black hole evaporation might lead to different numbers of particles and antiparticles in outer world belongs to Hawking [13], while Zeldovich [14] proposed a concrete mechanism of its realization. Possible generation of the cosmological baryon asymmetry in the process of primordial black hole (PBH) evaporation and an estimate of its magnitude are described in the subsequent publications [15, 16]. The calculations of the baryon asymmetry of the universe in the frameworks of the Zeldovich scenario [14] have been performed in refs. [17, 18].

These pioneering works were followed by now with a plethora of scenarios according to which cosmological baryon asymmetry could be generated in the process of black hole evaporation [19–29]. In contrast to them the model presented in this paper is based on the novel idea that baryogenesis could proceed in the process of asymmetric capture of baryons and antibaryons by primordial black holes. Here the ideas suggested in our previous paper [30] are further developed. In ref. [30] we considered slow diffusion of heavy baryons and antibaryons to primordial black holes because of their short mean free
path in the primeval plasma. We have shown that a non-zero value of the baryon asymmetry may be generated in this process, though a strong fine-tuning was necessary to get a reasonable result.

Here we consider the scenario with a large mean free path of the accreting baryons and because of that the model becomes much less constrained allowing for generation of the observed cosmological baryon asymmetry.

The main point of the present work is the difference of mobilities of heavy nonrelativistic baryons in the primeval plasma predominantly consisting of relativistic matter. This difference can be induced by the breaking of C and CP invariance, though CPT remains unbroken in the same way as the partial decay width of particles and antiparticles may be different if a sufficient number of the decay channels are open, see e.g. [31]. A detailed discussion of the cross-section difference is presented in our paper [30]. Hence the accretion rate becomes different and antibaryons can accumulate inside PBHs leaving excess of baryons in our external space.

The paper is organized as follows. In the following Section we present and solve the equation of motion of nonrelativistic particles accreting to a central gravitating body (BH) in expanding Friedmann universe. In Sec. II cosmological baryon asymmetry due to antibaryon capture by PBH is calculated and the values of the parameters are fixed to ensure the proper magnitude of the asymmetry. Throughout this paper these particles are called X and ¯X for baryons and antibaryons respectively. In Sec. IV we describe essential features of C and CP symmetries violation necessary for the implementation of the considered scenario (for more detailed discussion see ref. [30]). In Sec. V we conclude.

II. ACCRETION TO BH IN EXPANDING UNIVERSE

Equation of motion (geodesic equation) for nonrelativistic particles, X, in the curved background created by a black hole in the Friedmann space-time is derived e.g. in ref. [32] and has the form:

\[ \ddot{r} = \frac{\ddot{a}}{a} - \frac{r_g}{2a^2} + \frac{L^2}{r^3}, \]

where \( r_g \) is the Schwarzschild radius, \( r_g = 2MG_N = 2M/m_{pl}^2 \), \( M \) is the PBH mass, \( G_N \) is the Newtonian gravitational constant, \( m_{pl} = 1.22 \cdot 10^{19} \text{ GeV} = 2.17 \cdot 10^{-5} \text{ g} \), \( L \) is the angular momentum the X-particle, \( a(t) \) is the cosmological scale factor, and the over-dots mean time derivatives. At the radiation dominated cosmological stage \( \ddot{a}/a \approx -H^2 \) if the PBH does not have an essential impact on the Friedmann expansion. The angular momentum of the X-particle, \( L \), is supposed to be zero since these particles and PBH are naturally at rest in the comoving frame.

We integrate this equation analytically assuming that \( H \) slowly, adiabatically changes with time. As a result we find the following expression for the X-particle velocity, \( \dot{r} \):

\[ r^2 = \frac{r_g}{r} - H^2 r^2 + v_1^2, \]

where \( v_1 = \text{const} \). We fix this constant by the condition that the particle velocity vanishes at the distance \( r = r_{\text{max}} \) corresponding to the equilibrium between the Hubble repulsion and the gravitational attraction induced by the BH so \( r_{\text{max}} \) is the maximum radius of particle capture. As follows from eq. (3), \( r_{\text{max}} \) is equal to

\[ r_{\text{max}}^3 = \frac{r_g}{H^2}. \]

So \( v_1 = 0 \) and eq. (3) is solved as

\[ \dot{r} = -\sqrt{\frac{r_g}{r} - H^2 r^2}. \]

Equation (6) can be further integrated resulting in:

\[ r(t) = r_{\text{max}} \left[ \cos(3Ht/2) \right]^{2/3}. \]

It is valid till \( r(t) \) drops down to \( r_g \). Since \( r_g \ll r_{\text{max}} \), it can be reached at \( 3Ht/2 \) close to \( \pi/2 \). It does not contradict our assumption of slow variation of \( H \) which is true by an order of magnitude if \( Ht \lesssim 1 \)

In the free fall approximation there is no difference between the laws of motion for \( X \) and \( \bar{X} \). To take their difference into account we need to include into equation of motion (2) a small friction term induced by the interaction of \( X \) and \( \bar{X} \) particles with the plasma surrounding a black hole. With this correction the trajectory of \( X \)-particles, \( \bar{R}(t), \) would obey the equation:

\[ \ddot{\bar{R}} = \frac{\ddot{a}}{a} \bar{R} - \frac{r_g}{2R^3} - \gamma \dot{\bar{R}}, \]

(7)
where \( \gamma = \sigma_{el} v n_{rel} \), \( \sigma_{el} \sim f^4/m_X^3 \) is the cross-section of elastic \( X \)-particle scattering by the relativistic \( X \) ones in cosmic plasma, \( f \) is the coupling constant, \( v \approx 1 \) is the relative velocity of \( X \) and a relativistic scatterer, and \( n_{rel} = 0.1 g_s T^3 \) is the number density of all relativistic species in plasma, \( g_s \approx 100 \) is the number of relativistic species, and \( T \) is the plasma temperature. Due to C and CP violation \( \gamma \) should be different for \( X \) and \( \bar{X} \), i.e. \( \Delta \gamma = \gamma_X - \gamma_{\bar{X}} \neq 0 \).

Assuming that \( \gamma/H \) is small, we solve eq. (7) perturbatively expanding \( R \) as \( R = r + r_1 \) where \( r \) is the solution of the equation of motion (7) in zeroth order in \( \gamma \), where eq. (9) does not change this result significantly. Thus \( \gamma/H \) appears to be smaller than the second one. Hence this equation can be solved as:

where \( T < m_X \) very small temperatures, \( \Delta \eta = \eta_X - \eta_{\bar{X}} \neq 0 \).

The total number of \( X \) or \( \bar{X} \) particles captured by a BH during the Hubble time is approximately:

\[
N \approx (4\pi/3) (1 + 3\gamma t) r_{max}^3 n_X, \tag{12}
\]

where \( n_X \) is the number density of \( X \) particles. If the annihilation \( X \bar{X} \) is weak (we check below when it is indeed the case) and if \( X \)-particles are efficiently produced by the inflaton decay at the end of inflation, then

\[
n_X \approx n_{\bar{X}} \approx 0.1 g_s T^3, \tag{13}
\]

where \( g_s \) is the number of spin states of \( X \)-particles.

Since the \( X \bar{X} \)-annihilation is weak, the number density of \( X \)-particles remains unsuppressed even at very small temperatures, \( T < m_X \).

III. BARYOGENESIS THROUGH CAPTURE OF BARYONS BY PBHS

We assume that the heavy particles \( X \) and antiparticles \( \bar{X} \) have non-zero baryon number \( B_X \sim 1 \). We also assume that there exists an interaction between \( X \), \( \bar{X} \), and light particles which breaks C and CP symmetries but respects CPT. Some other particles, heavier than \( X \) and rather short-lived are also needed. Their existence is necessary to create difference between elastic cross-sections of \( X \) and \( \bar{X} \) particles in the primeval plasma despite the CPT restrictions which demand equality of the total cross-sections, see Sec. [IV]

The accretion of \( X \)-particles to a PBH effectively started when these particles became nonrelativistic. As we see in what follows, the smaller the plasma temperature, the larger the baryon asymmetry, if the density of \( X \)-particles in comoving volume does not drop as \( \exp(-m_X/T) \). In other words it happens if the \( X \bar{X} \)-annihilation froze at temperatures of the order of \( m_X \). However, such early freezing of massive species, if they are stable, would create too high contribution into the cosmological density of dark matter. The problem can be solved if \( X \)-particles are unstable but live sufficiently long to fulfill their task of creating the baryon asymmetry of the universe. These and some other constraints are considered below in this section.

Using equation (12) we can estimate the difference between number of captured \( X \) and \( \bar{X} \) particles by a single PBH during time interval \( t \):

\[
\Delta N \approx 4\pi r_{max}^3 n_X t \Delta \gamma, \tag{14}
\]
where \( n_X \) is the number density of \( X \)-particles after they became nonrelativistic. We should keep in mind that this time duration is bounded by the condition \( t \lesssim H(T) \), where \( T \) is the favorable temperature for excessive \( X \) over \( \bar{X} \) capture by PBH, see below.

The difference between the friction coefficients in the case of maximally broken C and CP symmetries can be estimated as

\[
\Delta \gamma = \delta \sigma_{el} n_{rel} \approx f^6 n_{rel} / m_X^2 
\]

since the cross section of elastic scattering of \( X \)-particles on the relativistic particles is \( \sigma_{el} \approx f^4 / m_X^2 \) and the difference between \( X \) and \( \bar{X} \) scattering is of the order of \( f^2 \) because the cross-section difference appears as a result of radiative correction proportional to \( f^2 \), exactly as there appears the difference between partial decay widths in the scenario of baryogenesis through massive particle decays.

This result is true if the following conditions are fulfilled: the mean free path of \( X \)-particles in the primordial plasma \( l_{\text{free}} \) should be larger than the maximum capture radius \( r_{\text{max}} \). The former can be estimated as:

\[
l_{\text{free}} = \frac{1}{\sigma_{el} n_{rel}} = \frac{m_X^2}{0.1 g_* T^3 f^2} 
\]

where \( \sigma_{el} \) and \( n_{rel} \) are defined below eq. (7).

The condition \( l_{\text{free}} > r_{\text{max}} \) can be rewritten as

\[
\frac{m_X^2 H}{0.1 g_* f^2 T^3} > (r_H H)^{1/3} 
\]

The Hubble parameter at the expansion stage dominated by relativistic matter can be expressed through the temperature applying the set of the following equations:

\[
\varrho = \frac{3 H^2 m_{Pl}^2}{8 \pi} = \frac{3 m_{Pl}^2}{32 \pi t^2} = \frac{\pi^2 g_*}{30} T^4. 
\]

Hence we obtain:

\[
H = \left( \frac{8 \pi^3 g_*}{90} \right)^{1/2} T^2 = 16.6 g_{100}^{1/2} T^2 m_{Pl}^2, 
\]

where \( g_{100} = g_*/100 \).

Thus the bound (17) can be rewritten as:

\[
M < \frac{0.14 m_X^6}{g_{100}^2 f^{12} T^6} 
\]

Note that the limit does not depend upon the Planck mass.

If \( m_X = 3 \times 10^{13} \text{ GeV} \) (close to the typical heating temperature after inflation), \( T = m_X \), and \( f = 0.1 \), the condition of the free fall is fulfilled for \( M \lesssim 7 \text{ g} \). For higher mass of the PBHs the free fall condition is satisfied at smaller \( T \), e.g. if \( M = 10^6 \text{ g} \), the efficient free fall capture took place at \( T = m_X/10 \).

The rate of the annihilation is determined by the equation:

\[
\Gamma_{\text{ann}} \equiv \dot{n}_X / n_X = \sigma_{\text{ann}} v n_X = 0.1 g_* f_{\text{ann}}^4 T^3 / m_X^2 
\]

where \( f_{\text{ann}} \) is the coupling constant of the annihilation. Demanding that \( \Gamma_{\text{ann}} \) is small in comparison with \( H \), see eq. (19), we find that the annihilation would be inefficient at the temperatures satisfying the condition:

\[
\frac{T}{m_X} < 2.5 \times 10^{-6} f_{\text{ann}}^{-4} \left( \frac{m_X}{3 \times 10^{13} \text{ GeV}} \right). 
\]

If the maximum value of \( T/m_X \) may reach unity, then the annihilation does not essentially diminish the density of \( X \)-particles below \( n_X = 0.1 g_* T^3 \). In other words the density of \( X \) and \( \bar{X} \) particles would be conserved in the comoving volume, i.e. \( n_X = 0.1 g_* T^3 \) below \( T = m_X \). For \( m_X = 3 \times 10^{13} \text{ GeV} \) this could be realized if \( f_{\text{ann}} \lesssim 4 \times 10^{-2} \). Otherwise the density of \( X \)-particles would be exponentially suppressed, \( n_X \sim \exp(-m_X/T) \) at low temperatures. To avoid an overclosing of the universe by \( X \)-particles we assume that they are unstable, presumably decaying before the Big Bang Nucleosynthesis.
Another important restriction on the efficiency of the discussed here mechanism is that the "size" of X-particles i.e. its Compton wave length $\lambda_X$ should be smaller than the gravitational radius of PBH. Otherwise the probability of the particle capture would be suppressed by a power of the ratio $r_g/\lambda_X$:

$$\lambda_X = 1/m_X < r_g = 2M/m_{Pl}^2.$$  

It leads to a lower bound on the PBH mass:

$$M > m_{Pl}^2/2m_X = 4.4g \left(\frac{3 \cdot 10^{13}\text{GeV}}{m_X}\right).$$  

The baryon asymmetry gained by the PBH antibaryon capture can be diluted by the entropy release from the PBH evaporation. As it follows from Ref. [33] this would be avoided if

$$\epsilon M < 10^{-5} g,$$

where $\epsilon$ is the fraction of the energy density of PBHs at the moment of their formation:

$$\frac{\rho_{PBH}(t_{\text{form}})}{\rho_{\text{rel}}(t_{\text{form}})} = \epsilon,$$

where $\rho_{\text{rel}} \approx 3Tn_{\text{rel}}$ is the energy density of the relativistic matter, and

$$t_{\text{form}} = M/m_{Pl}^2.$$  

Using eqs. (18) and (19) we find that the temperature of the relativistic matter at the formation moment is

$$T_{\text{form}} = T(t_{\text{form}}) = 0.17g_{100}^{-1/4}m_{Pl}\left(\frac{m_{Pl}}{M}\right)^{1/2} = 10^{14}\text{GeV}g_{100}^{-1/4}M^{-1/2},$$

where $M = M/10^4$. For successful baryogenesis PBHs should be created while X-particles are abundant in the cosmological plasma. If $X\bar{X}$-annihilation continued till $T \ll m_X$, then the temperature of the PBH creation should be not much smaller than $m_X$. If, as we assume in the present work, the annihilation of $X\bar{X}$ in thermal plasma was never efficient, then PBHs should be produced prior to the decay of $X$ (and $X$)-particles.

In the course of the cosmological expansion and cooling down the energy fraction of PBH rises as ($T_{\text{form}}/T$) until their evaporation, which happens at the time moment equal to the BH life-time $\tau_{BH}$ [33]:

$$t = \tau_{BH} = 30M^3/m_{Pl}^4.$$  

The corresponding temperature is

$$T(\tau_{BH}) = 0.03g_{100}^{-1/4}m_{Pl}\left(\frac{m_{Pl}}{M}\right)^{3/2} = 3.7 \cdot 10^4\text{GeV}M^{-3/2}.$$  

The temperature of the relativistic plasma at the moment of PBH decay should be smaller than $m_X$ to allow nonrelativistic X-particles to be captured by PBHs.

Since PBH are nonrelativistic, while the bulk of the matter is relativistic, the fraction of the PBH energy density at temperature $T$ becomes larger than $\epsilon$ by the factor

$$\frac{T_{\text{form}}}{T} \approx g_{100}^{-1/4}M^{-1/2}10^{14}\text{GeV}.$$  

Now using Eq. (14) and Eq. (15) we find for the excessive baryon number create by a single primordial black holes:

$$\Delta N = 4\pi f^6\frac{r_g n_{\text{rel}} n_X t}{H^2 m_X^2} = \frac{9g_{s}f^6}{\pi^2} MT^2 t.$$  

Hence the baryon asymmetry can be estimated as

$$\beta = \frac{B_{Xn_{BH}}\Delta N}{n_{\text{rel}}} = \frac{2.7}{\pi^2} B_Xg_sf^6\frac{\rho_{PBH}}{\rho_{\text{rel}}}T^3 t = \frac{2.7}{\pi^2} B_Xg_sf^6\epsilon T_{\text{form}}T^2 t$$

where $\epsilon$ is the fraction of PBH energy density to that of the relativistic matter at the moment of their formation, see eq. (26). $T_{\text{form}}$ is the temperature at PBH formation (28), and and $t \sim 1/H$, so we finally obtain:

$$\beta \approx 0.016B_Xg_sf^6\epsilon \frac{T_{\text{form}} m_{Pl} T}{T m_X^2}.$$  

Taking the maximum allowed values of $\epsilon$ from Eq. (25) $\epsilon = 10^{-5} g/M$, $m_X = 10^{13}$ GeV, $f = 0.1$, $T = m_X/10$ $M = 10^4$ g and thus $T_{\text{form}} = 10^{14}$ GeV we find that the baryon asymmetry can easily reach the observed value and even overcome it. This choice of the parameters satisfies the derived above restrictions.
among the partial modes of scattering processes, but the magnitude of the latter would be suppressed and antiparticles while allows for a difference between the partial decay rates as well as for a difference \[ \lambda \]

As \( X \) scattering of \( \bar{X} \) scattering between the partial decay widths of particles and antiparticles, we consider radiative corrections to elastic interaction leading to disappearance of \( \bar{X} \) particles in cosmic plasma. The complex coefficient \( \Gamma \) where summations are done over all light particle sets in the initial state

where\( \sum \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \)

where \( \sum \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \) \( \bar{X} \) \( \rightarrow \) \( \bar{X} \)

In complete analogy with the higher order corrections to the decay process, which create a difference between the probabilities of charge conjugated processes can be estimated as:

where \( \sigma_X \) and \( \sigma_{\bar{X}} \) stand respectively for the scattering cross-section of \( X \) and \( \bar{X} \) particles on relativistic particles, which lead to different values of the cross-sections. For more detail see our earlier paper \[ 30 \]. The corresponding Feynman diagrams are presented in Figs. 1(a) and 1(b), an example of the process leading to

\[ a + X \rightarrow b + Y, \]

where the heavy particle \( Y \) may have zero baryonic number and the light state \( b \) should have the same baryonic number as \( a + X \), if we want to avoid non-conservation of baryons.

In complete analogy with the higher order corrections to the decay process, which create a difference between the partial decay widths of particles and antiparticles, we consider radiative corrections to elastic scattering of \( X \) on relativistic particles, which lead to different values of the cross-sections. For more detail see our earlier paper \[ 30 \]. The corresponding Feynman diagrams are presented in Figs. 1(a) and 1(b), an example of the process leading to

\[ \text{FIG. 1: Feynman diagrams describing } X \text{ (or } \bar{X} \text{) scattering off light quark. The first, 'a', diagram is the lowest order contribution and gives equal values for particles and antiparticles. The 'b' diagram present an example of the one-loop correction with an exchange of } Y \text{-particle. The one-loop scattering may be also present to the initial state and its contribution multiplies the result by factor } 2. \]

The equalities of the total probabilities means in particular that the total cross-section of \( X \)-scattering on particle "a" \( \sigma_{tot}(X + a \rightarrow All) \) is equal to the same of particles and antiparticles \( \sigma_{tot}(\bar{X} + \bar{a} \rightarrow All) \). If the final state "All" contains one and only one \( X \)-particle, then the mobilities of \( X \) and \( \bar{X} \) in the cosmological plasma would be the same and the discussed here mechanism of baryogenesis would not operate. However, if the complete set of the final states includes a state or states where \( X \) is missing (and analogously for the reactions with \( \bar{X} \)), then the cross-sections of the processes \( \sigma_{tot}(X + a \rightarrow X + All) \) and \( \sigma_{tot}(\bar{X} + \bar{a} \rightarrow \bar{X} + All) \) may be different, leading to the needed mobility differences of \( X \) and \( \bar{X} \).

The difference between probabilities of charge conjugated processes can be estimated as:

\[ \delta = \frac{\sigma_X X_0 - \sigma_{\bar{X}} \bar{X}_0}{\sigma_X X_0 + \sigma_{\bar{X}} \bar{X}_0} \approx \frac{|g_{x1}|^2 \text{Im}(D) \text{Im}(g_{y1} g_{y2}^2 g_{y1} g_{y2})}{|g_{x1}|^4} \propto \alpha f^2, \]

where \( \sigma_X \) and \( \sigma_{\bar{X}} \) stand respectively for the scattering cross-section of \( X \) and \( \bar{X} \) particles on relativistic particles in cosmic plasma. The complex coefficient \( D \) comes from the integration over the loop, and \( g_{x1} \) and \( g_{y1} \) are partial decay constants of \( X \) and \( Y \) particles respectively.
The following supersymmetry inspired model can serve as appropriate frameworks for the scenario. Assume that $X$ is an analogue of the lightest supersymmetric particle (LSP) which is stable due to an analogue of $R$-parity. Let assume that there exists a heavier partner $H$ with zero baryonic number which would be unstable and decay through the channel $H \to X + 3q$, where $q$ are light quarks with proper quantum numbers. Accordingly the reaction $X + q \to H + 2q$ becomes possible. It is exactly what we need to allow for a difference between the cross-sections of the reaction $\sigma_{tot}(X + a \to X + All)$ and $\sigma_{tot}(\bar{X} + \bar{a} \to \bar{X} + All)$, which can lead to a different mobilities of $X$ and $\bar{X}$ around black hole and to dominant capture of antibaryons over baryons creating cosmological baryon asymmetry.

Note that in $R^2$ gravity, Starobinsky inflation \cite{36}, the allowed mass of LSP-kind particle can be close to $10^{13}$ GeV or even higher \cite{37,38}.

V. CONCLUSION

In this paper we continue investigation of baryogenesis through the asymmetric capture of baryons and antibaryons by primordial black holes. Unlike in our previous paper \cite{30}, where we used diffusion approximation $\gamma/H \gg 1$ in which particles many times scatter off the cosmic plasma before they are captured by PBH, in this study we investigate the opposite limit in which $\gamma/H \ll 1$ or the free fall limit. As it appears, there is a sufficiently wide parameters space to explain the observed value of baryon asymmetry of the universe.

A noticeable increase of the baryon asymmetry generated by the capture of the antibaryonic number by PBH in the considered version of the scenario is achieved due to assumed negligible annihilation of $X$-particles with decreasing temperatures, $T < m_X$, because at smaller $T$ the relative fraction of PBH with respect to the total cosmological energy density goes up quite significantly.

The proposed here mechanism of baryogenesis does not demand two out of three Sakharov’s principles. Namely it can proceed in thermal equilibrium and without assumption of non-conservation of baryonic number in particle reactions. It helps to avoid a possible problem which arises because non-conservation of baryons is not (yet) observed in direct experiment.

In a sense black holes break conservation of baryonic number, either hiding baryons in internal space making them unobservable, if black holes are eternal, or transforming an arbitrary amount of baryons into a state with zero baryonic number. For instance a black hole consisting entirely from baryons would completely evaporate creating (almost) equal number of baryons and antibaryons. In the process of evaporation a small baryon asymmetry might be created but it normally would be negligibly small in comparison with the initial baryonic number captured at the black hole formation.

If baryonic number is conserved in particle interactions, the proton must be almost absolutely stable. To be more precise it may decay by Zeldovich mechanism \cite{10} through formation of a virtual black hole from three quarks inside proton. But the life-time with respect to such decay is almost infinite, $\tau_p \sim 10^{45}$ years. Also one could hardly expect neutron-antineutron oscillations induced by virtual BHs to be observable (for a recent review see e.g. \cite{39}).

Another unusual feature of the model is a possibility to create baryon (or any other type of asymmetry between particles and antiparticles) in thermal equilibrium. Normally the deviation from thermal equilibrium is suppressed by the factor of the order of the ratio of the Hubble expansion rate to the particle reaction rate, $H/T$. The former is inversely proportional to a huge value of the Planck mass, $H \sim T^2/m_{Pl}$, where $T$ is the cosmological plasma temperature. According to the estimates presented above, for the mechanism considered here the situation is opposite: the larger is the Planck mass (or the slower is the cosmological expansion), the larger is the baryon asymmetry. On the other hand, the gravitational attraction which forces massive $X$-particles to fall on the nearest BH is inversely proportional to $m_{Pl}^2$, so ultimately the effect disappears in the limit of infinite $m_{Pl}$ as well.

The magnitude of the baryon asymmetry evidently strongly depends upon the cosmological expansion regime. In particular, it would be very interesting to study baryogenesis in the frameworks of $R^2$ inflation \cite{34} there exists a long period of the universe evolution during which the fall off of the cosmological temperature is drastically different from that accepted in the conventional cosmology \cite{37,38}.

In the course of working on the presented here version of baryogenesis we became aware of an interesting modification of the scenario presented in ref. \cite{30} on the generation of the cosmological baryon asymmetry through the capture of antibaryons by PBH \cite{41}, which also may lead to an efficient baryogenesis.

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