Spontaneous and Gravitational Baryogenesis

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Some problems of spontaneous and gravitational baryogenesis are discussed. Gravity modification due to the curvature dependent term in gravitational baryogenesis scenario is considered. It is shown that the interaction of baryonic fields with the curvature scalar leads to strong instability of the gravitational equations of motion and as a result to noticeable distortion of the standard cosmology.

Keywords: cosmology, baryogenesis, gravitational equations, modified theories of gravity.

1. Introduction

Observations show that at least the region of the Universe around us is matter-dominated. Though we understand how the matter-antimatter asymmetry may be created, the concrete mechanism is yet unknown. The amount of antimatter is very small and it can be explained as the result of high energy collisions in space. The existence of large regions of antimatter in our neighbourhood would produce high energy radiation as a consequence of matter-antimatter annihilation, which is not observed. Any initial asymmetry at inflation could not solve the problem of observed excess of matter over antimatter, because the energy density associated with baryonic number would not allow for sufficiently long inflation.

On the other hand, matter and antimatter seem to have similar properties and therefore we could expect a matter-antimatter symmetric Universe. A satisfactory model of our Universe should be able to explain the origin of the local observed matter-antimatter asymmetry. The term baryogenesis means the generation of the asymmetry between baryons (basically protons and neutrons) and antibaryons (antiprotons and antineutrons).

In 1967 Andrey Sakharov pointed out 3 ingredients, today known as Sakharov principles, to produce a matter-antimatter asymmetry from an initially symmetric Universe. These conditions include: 1) non-conservation of baryonic number; 2) breaking of symmetry between particles and antiparticles; 3) deviation from thermal equilibrium. However, not all of three Sakharov principles are strictly necessary.

In what follows we briefly discuss some features of spontaneous baryogenesis (SBG) and concentrate in more detail on gravitational baryogenesis (GBG). Both these mechanisms do not demand an explicit C and CP violation and can proceed in thermal equilibrium. Moreover, they are usually most efficient in thermal
The statement that the cosmological baryon asymmetry can be created by spontaneous baryogenesis in thermal equilibrium was mentioned in the original paper by Cohen and Kaplan \(^1\) and developed in subsequent papers\(^2,3\), for review see \(^4,5\).

The term "spontaneous" is related to spontaneous breaking of a global \(U(1)\)-symmetry, which ensures the conservation of the total baryonic number in the unbroken phase. This symmetry is supposed to be spontaneously broken and in the broken phase the Lagrangian density acquires the additional term

\[
\mathcal{L}_{SB} = (\partial_\mu \theta) J^\mu_B ,
\]

where \(\theta\) is the Goldstone field and \(J^\mu_B\) is the baryonic current of matter fields, which becomes non-conserved.

For a spatially homogeneous field, \(\theta = \theta(t)\), the Lagrangian is reduced to the simple form

\[
\mathcal{L}_{SB} = \dot{\theta} n_B , \quad n_B \equiv J^0_B ,
\]

where time component of a current is the baryonic number density of matter, so it is tempting to identify \(\dot{\theta}\) with the chemical potential, \(\mu_B\), of the corresponding system. However, such identification is questionable and depends upon the representation chosen for the fermionic fields\(^6,7\). It is heavily based on the assumption \(\dot{\theta} \approx \text{const}\), which is relaxed in the work\(^8\). But still the scenario is operative and presents a beautiful possibility to create an excess of particles over antiparticles in the Universe.

Subsequently the idea of gravitational baryogenesis (GBG) was put forward\(^9\), where the scenario of SBG was modified by the introduction of the coupling of the baryonic current to the derivative of the curvature scalar \(R\):

\[
\mathcal{L}_{GBG} = \frac{1}{M^2} (\partial_\mu R) J^\mu_B ,
\]

where \(M\) is a constant parameter with the dimension of mass.

In the presented talk we demonstrate that the addition of the curvature dependent term \(^3\) to the Hilbert-Einstein Lagrangian of General Relativity (GR) leads to higher order gravitational equations of motion, which are strongly unstable with respect to small perturbations. The effects of this instability may drastically distort not only the usual cosmological history, but also the standard Newtonian gravitational dynamics. We discovered such instability for scalar baryons\(^10\) and found similar effect for the more usual spin one-half baryons (quarks)\(^11\).

2. Gravitational baryogenesis with scalar baryons

Let us start from the model where baryonic number is carried by scalar field \(\phi\) with potential \(U(\phi, \phi^*)\). An example with baryonic current of fermions will be considered in the next section.

The action of the scalar model has the form:

\[
A = \int d^4x \sqrt{-g} \left[ \frac{m^2}{16\pi} R + \frac{1}{M^2} (\partial_\mu R) J^\mu_B - g^{\mu\nu} \partial_\mu \phi \partial_\nu \phi^* + U(\phi, \phi^*) \right] - A_m ,
\]
where \( m_{Pl} = 1.22 \cdot 10^{19} \) GeV is the Planck mass, \( A_m \) is the matter action, \( J^\mu = g^{\mu\nu} J_\nu \), and \( g^{\mu\nu} \) is the metric tensor of the background space-time. We assume that initially the metric has the usual GR form and study the emergence of the corrections due to the instability described below.

In contrast to scalar electrodynamics, the baryonic current of scalars is not uniquely defined. In electrodynamics the form of the electric current is dictated by the conditions of gauge invariance and current conservation, which demand the addition to the current of the so called sea-gull term proportional to \( e^2 A_\mu |\phi|^2 \), where \( A_\mu \) is the electromagnetic potential.

On the other hand, a local \( U(1) \)-symmetry is not imposed on the theory determined by action (4). It is invariant only with respect to a \( U(1) \) transformations with a constant phase. As a result, the baryonic current of scalars is considerably less restricted. In particular, we can add to the current an analog of the sea-gull term, \( \sim (\partial_\mu R) |\phi|^2 \), with an arbitrary coefficient.

In our paper\(^{10}\) we study the following two extreme possibilities, when the sea-gull term is absent and the current is not conserved, or the sea-gull term is included with the coefficient ensuring current conservation. In both cases no baryon asymmetry can be generated without additional interactions. It is trivially true in the second case, when the current is conserved, but it is also true in the first case despite the current non-conservation, simply because the non-zero divergence \( D_\mu J^\mu \) does not change the baryonic number of \( \phi \) but only leads to redistribution of particles \( \phi \) in the phase space. So to create any non-zero baryon asymmetry we have to introduce an interaction of \( \phi \) with other particles which breaks conservation of \( B \) by making the potential \( U \) non-invariant with respect to the phase rotations of \( \phi \), as it is described below.

If the potential \( U(\phi) \) is not invariant with respect to the \( U(1) \)-rotation, \( \phi \to \exp \left(i\beta\right) \phi \), the baryonic current defined in the usual way

\[
J^\mu = iq (\partial^\mu \phi - \phi \partial^\mu \phi^* )
\]

(5)

is not conserved. Here \( q \) is the baryonic number of \( \phi \) and we omitted index \( B \) in current \( J^\mu \).

With this current and Lagrangian \(^{11}\) the equation for the curvature scalar, \( R \), takes the form:

\[
\frac{m_{Pl}^2}{16\pi} R + \frac{1}{M^2} \left[ (R + 3D^2) D_\alpha J^\alpha + J^\alpha D_\alpha R \right] - D_\alpha \phi D^\alpha \phi^* + 2U(\phi) = -\frac{1}{2} T^\mu_\mu ,
\]

(6)

where \( D_\mu \) is the covariant derivative in metric \( g_{\mu\nu} \) (of course, for scalars \( D_\mu = \partial_\mu \)) and \( T_\mu^\nu \) is the energy-momentum tensor of matter obtained from action \( A_m \).

According to definition (5), the current divergence is:

\[
D_\mu J^\mu = \frac{2q^2}{M^2} \left[ D_\mu R (\partial^\mu \phi + \phi \partial^\mu \phi^*) + |\phi|^2 D^2 R \right] + iq \left( \phi \frac{\partial U}{\partial \phi} - \phi^* \frac{\partial U}{\partial \phi^*} \right).
\]

(7)

If the potential \( U \) is invariant with respect to the phase rotation of \( \phi \), i.e. \( U = U(|\phi|) \), the last term in this expression disappears. Still the current remains
non-conserved, but this non-conservation does not lead to any cosmological baryon asymmetry. Indeed, the current non-conservation is proportional to the product $\phi^*\phi$, so it can produce or annihilate an equal number of baryons and antibaryons.

To create cosmological baryon asymmetry we need to introduce new types of interactions, for example, the term in the potential of the form: $U_4 = \lambda_4\phi^4 + \lambda_4^*\phi^4$. This potential is surely non invariant w.r.t. the phase rotation of $\phi$ and can induce the B-non-conserving process of transition of two scalar baryons into two antibaryons, $2\phi \rightarrow 2\bar{\phi}$.

Let us consider solution of the above equation of motion in cosmology. The metric of the spatially flat cosmological FRW background can be taken as:

$$ ds^2 = dt^2 - a^2(t)dx^2. $$

In the homogeneous case the equation for the curvature scalar (6) takes the form:

$$ m^2_{pl}\frac{8\pi R}{16\pi} + \frac{1}{M^2} \left[ (R + 3\dot{R}^2 + 9H\partial_t)D_\alpha J^\alpha + \dot{R} J^0 \right] = - \frac{T^{(tot)}}{2}, $$

where $J^0$ is the baryonic number density of the $\phi$-field, $H = \dot{a}/a$ is the Hubble parameter, and $T^{(tot)}$ is the trace of the energy-momentum tensor of matter including contribution from the $\phi$-field. In the homogeneous and isotropic cosmological plasma

$$ T^{(tot)} = \rho - 3P, $$

where $\rho$ and $P$ are respectively the energy density and the pressure of the plasma. For relativistic plasma $\rho = \pi^2g_*T^4/30$ with $T$ and $g_*$ being the plasma temperature and the number of particle species in the plasma. The Hubble parameter is expressed through $\rho$ as $H^2 = 8\pi\rho/(3m^2_{pl}) \sim T^4/m^2_{pl}$.

The covariant divergence of the current is given by the expression (11). In the homogeneous case we are considering it takes the form:

$$ D_\alpha J^\alpha = \frac{2q^2}{M^2} \left[ \dot{\bar{R}} (\phi^* \dot{\phi} + \phi \dot{\phi}^*) + (\bar{R} + 3H\dot{R}) \phi^* \phi \right] + iq \left( \phi \frac{\partial U}{\partial \phi} - \dot{\phi}^* \frac{\partial U}{\partial \phi^*} \right). $$

To derive the equation of motion for the classical field $R$ in the cosmological plasma we have to take the expectation values of the products of the quantum operators $\phi$, $\phi^*$, and their derivatives. Performing the thermal averaging, we find

$$ \langle \phi^* \phi \rangle = \frac{T^2}{12}, \quad \langle \phi^* \dot{\phi} + \dot{\phi}^* \phi \rangle = 0. $$

Substituting these average values into Eq. (11) and neglecting the last term in Eq. (11) we obtain the fourth order differential equation:

$$ m^2_{pl}\frac{8\pi R}{16\pi} + \frac{q^2}{6M^4} (R + 3\dot{R}^2 + 9H\partial_t) \left[ (\bar{R} + 3H\dot{R}) T^2 \right] + \frac{1}{M^2}\dot{R} \langle J^0 \rangle = - \frac{T^{(tot)}}{2}. $$

Here $\langle J^0 \rangle$ is the thermal average value of the baryonic number density of $\phi$. It is assumed to be zero initially and generated as a result of GBG. We neglect this term,
since it is surely small initially and probably subdominant later. Anyhow it does not noticeably change the exponential rise of $R$ at the onset of the instability.

Eq. (13) can be further simplified if the variation of $R(t)$ is much faster than the universe expansion rate or in other words $\dot{R}/R \gg H$. Correspondingly the temperature may be considered adiabatically constant. The validity of these assumptions is justified a posteriori after we find the solution for $R(t)$.

Keeping only the linear in $R$ terms and neglecting higher powers of $R$, such as $R^2$ or $HR$, we obtain the linear differential equation of the fourth order:

$$\frac{d^4R}{dt^4} + \mu^4 R = -\frac{1}{2} T^{(tot)}$$

where $\mu^4 = \frac{m_{Pl}^2 M^4}{8\pi q^2 T^2}$. (14)

The homogeneous part of this equation has exponential solutions $R \sim \exp(\lambda t)$ with

$$\lambda = |\mu| \exp(i\pi/4 + i\pi n/2),$$

where $n = 0, 1, 2, 3$.

There are two solutions with positive real parts of $\lambda$. This indicates that the curvature scalar is exponentially unstable with respect to small perturbations, so $R$ should rise exponentially fast with time and quickly oscillate around this rising function.

Now we need to check if the characteristic rate of the perturbation explosion is indeed much larger than the rate of the universe expansion, that is:

$$(Re \lambda)^4 > H^4 = \left(\frac{8\pi \rho}{3m_{Pl}^2}\right)^2 = \frac{16\pi^6 g^2 T^8}{2025 m_{Pl}^2}.$$

(16)

where $\rho = \pi^2 g_* T^4 / 30$ is the energy density of the primeval plasma at temperature $T$ and $g_* \sim 10 - 100$ is the number of relativistic degrees of freedom in the plasma. This condition is fulfilled if

$$\frac{2025}{2\pi^6 q^2 g_*^2} \frac{m_{Pl}^6 M^4}{T^{10}} > 1,$$

(17)

or, roughly speaking, if $T \lesssim m_{Pl}^3 / M^2$. Let us stress that at these temperatures the instability is quickly developed and the standard cosmology would be destroyed.

If we want to preserve the successful big bang nucleosynthesis (BBN) results and impose the condition that the development of the instability was longer than the Hubble time at the BBN epoch at $T \sim 1$ MeV, then $M$ should be extremely small, $M < 10^{-32}$ MeV. The desire to keep the standard cosmology at smaller $T$ would demand even tinier $M$. A tiny $M$ leads to a huge strength of coupling $\mathcal{O}$. It surely would lead to pronounced effects in stellar physics.
3. Gravitational baryogenesis with fermions

Let us now generalize results, obtained for scalar baryons, to realistic fermions. We start from the action in the form

$$A = \int d^4x \sqrt{-g} \left[ \frac{m_0^2}{16\pi} R - \mathcal{L}_m \right]$$

(18)

with

$$\mathcal{L}_m = \frac{i}{2} \left( \bar{Q} \gamma^\mu \nabla_\mu Q - \nabla_\mu \bar{Q} \gamma^\mu Q \right) - m_Q \bar{Q} Q$$

$$+ \frac{i}{2} \left( \bar{L} \gamma^\mu \nabla_\mu L - \nabla_\mu \bar{L} \gamma^\mu L \right) - m_L \bar{L} L$$

$$+ \frac{g}{m_X^2} \left[ (\bar{Q} Q^c) (\bar{Q} L) + (\bar{Q}^c Q) (\bar{L} Q) \right] + \frac{f}{m_0^2} (\partial_\mu R) J^\mu + \mathcal{L}_{other},$$

(19)

where $Q$ is the quark (or quark-like) field with non-zero baryonic number, $L$ is another fermionic field (lepton), $\nabla_\mu$ is the covariant derivative of Dirac fermion in tetrad formalism. $m_0$ is a constant parameter with dimension of mass and $f$ is dimensionless coupling constant which is introduced to allow for an arbitrary sign of the curvature dependent term in the above expression. $J^\mu = \bar{Q} \gamma^\mu Q$ is the quark current with $\gamma^\mu$ being the curved space gamma-matrices, $\mathcal{L}_{other}$ describes all other forms of matter. The four-fermion interaction between quarks and leptons is introduced to ensure the necessary non-conservation of the baryon number with $m_X$ being a constant parameter with dimension of mass and $g$ being a dimensionless coupling constant. In grand unified theories $m_X$ may be of the order of $10^{14} - 10^{15}$ GeV.

Varying the action (18) over metric, $g^{\mu\nu}$, and taking trace with respect to $\mu$ and $\nu$, we obtain the following equation of motion for the curvature scalar:

$$-\frac{m_0^2}{8\pi} R = m_Q \bar{Q} Q + m_L \bar{L} L + \frac{2g}{m_X^2} \left[ (\bar{Q} Q^c) (\bar{Q} L) + (\bar{Q}^c Q) (\bar{L} Q) \right]$$

$$- \frac{2f}{m_0^2} (R + 3D^2) D_\alpha J^\alpha + T_{other},$$

(20)

where $T_{other}$ is the trace of the energy momentum tensor of all other fields. At relativistic stage, when masses are negligible, we can take $T_{matter} = 0$. The average expectation value of the interaction term proportional to $g$ is also small, so the contribution of all matter fields may be neglected.

As we see in what follows, kinetic equation leads to an explicit dependence on $R$ of the current divergence, $D_\alpha J^\alpha$, if the current is not conserved. As a result we obtain 4th order equation for $R$.

As previously, we study solutions of Eq. (20) in cosmology in homogeneous and isotropic FRW background with the metric $ds^2 = dt^2 - a^2(t) dr^2$. The curvature is a function only of time and has only time component, has the form:

$$D_\alpha V^\alpha = (\partial_t + 3H) V^t,$$
where \( H = \dot{a}/a \) is the Hubble parameter.

As an example let us consider the reaction \( q_1 + q_2 \leftrightarrow \bar{q}_3 + l_4 \), where \( q_1 \) and \( q_2 \) are quarks with momenta \( q_1 \) and \( q_2 \), while \( \bar{q}_3 \) and \( l_4 \) are antiquark and lepton with momenta \( q_3 \) and \( l_4 \). We use the same notations for the particle symbol and for the particle momentum. The kinetic equation for the variation of the baryonic number density \( n_B \equiv J^t \) through this reaction in the FRW background has the form:

\[
(\partial_t + 3H)n_B = I_B^{\text{coll}},
\]

where the collision integral for space and time independent interaction is equal to:

\[
I_B^{\text{coll}} = -3B_q(2\pi)^4 \int d\nu_{q_1,q_2} d\nu_{\bar{q}_3,l_4} \delta^4(q_1 + q_2 - q_3 - l_4) \left[ |A(q_1 + q_2 \rightarrow \bar{q}_3 + l_4)|^2 f_{q_1} f_{q_2} - |A(\bar{q}_3 + l_4 \rightarrow q_1 + q_2)|^2 f_{\bar{q}_3} f_{l_4} \right],
\]

where \( A(a \rightarrow b) \) is the amplitude of the transition from state \( a \) to state \( b \), \( B_q \) is the baryonic number of quark, \( f_a \) is the phase space distribution (the occupation number), and

\[
d\nu_{q_1,q_2} = \frac{d^3 q_1}{2E_{q_1}(2\pi)^3} \frac{d^3 q_2}{2E_{q_2}(2\pi)^3},
\]

where \( E_q = \sqrt{q^2 + m^2} \) is the energy of particle with three-momentum \( q \) and mass \( m \). The element of phase space of final particles, \( d\nu_{\bar{q}_3,l_4} \), is defined analogously.

We neglect the Fermi suppression factors and the effects of gravity in the collision integral. This is generally a good approximation.

The calculations are strongly simplified if quarks and leptons are in equilibrium with respect to elastic scattering and annihilation. In this case their distribution functions take the form

\[
f = \frac{1}{e^{(E/T - \xi)} + 1} \approx e^{-E/T + \xi},
\]

with \( \xi = \mu/T \) being dimensionless chemical potential, different for quarks, \( \xi_q \), and leptons, \( \xi_l \).

The assumption of kinetic equilibrium is well justified since it is usually enforced by very efficient elastic scattering. Equilibrium with respect to annihilation, say, into two channels: \( 2\gamma \) and \( 3\gamma \), implies the usual relation between chemical potentials of particles and antiparticles, \( \mu = -\mu \).

The baryonic number density is given by the expression:

\[
n_B = \int \frac{d^3 q}{2E_q(2\pi)^3} (f_q - f_{\bar{q}}) = \frac{gsB_q}{6} \left( \mu^2 + \frac{\xi^3}{\pi^2} \right) = \frac{gsB_qT^3}{6} \left( \xi + \frac{\xi^3}{\pi^2} \right),
\]

where \( T \) is the cosmological plasma temperature, \( gs \) and \( B_q \) are respectively the number of the spin states and the baryonic number of quarks.
We can use another representation of the quark field:

$$Q_2 = \exp(\frac{ifR}{m_0^2})Q$$

(27)

analogously to what is done in our paper\(^8\). Written in terms of \(Q_2\) Lagrangian \((20)\) would not contain terms proportional to \(f/m_0^2\), but dependence on such terms would reappear in the interaction term as:

$$\frac{2g_c m_x^2}{m_0^2} \left[ e^{-3ifR/m_0^2} (\bar{Q}_2 Q_2^c) (\bar{Q}_2 L) + e^{3ifR/m_0^2} (\bar{Q}_2^c Q_2) (\bar{L} Q_2) \right].$$

(28)

Nevertheless we obtain the same fourth order equation for the evolution of curvature, as for non-rotated field \(Q\).

Since the transition amplitudes, which enter the collision integral, are obtained by integration over time of the Lagrangian operator \((28)\), taken between the initial and final states, the energy conservation delta-function in Eq. \((23)\) would be modified due to time dependent factors \(\exp[\pm 3ifR(t)/m_0^2]\). In the simplest case, which is usually considered in gravitational (and spontaneous) baryogenesis, a slowly changing \(\dot{R}\) is taken, so we can approximate \(R(t) \approx \dot{R}(t) t\). In this case the energy is not conserved but the energy conservation condition is trivially modified, as

$$\delta \left[ E(q_1) + E(q_2) - E(l_3) - E(l_4) - 3f\dot{R}(t)/m_0^2 \right].$$

(29)

Thus the energy is non-conserved due to the action of the external field \(R(t)\). Delta-function \((29)\) is not precise, but the result is pretty close to it, if \(\dot{R}(t)\) changes very little during the effective time of the relevant reactions.

If the dimensionless chemical potentials \(\xi_q\) and \(\xi_l\), as well as \(f\dot{R}(t)/m_0^2/T\), are small, the collision integral can be written as:

$$I_{coll}^B \approx C_I g_s^2 T^8 m_x^4 \left[ \frac{3f\dot{R}(t)}{m_0^2 T} - 3\xi_q + \xi_l \right].$$

(30)

where \(C_I\) is a positive dimensionless constant. The factor \(T^8\) appears for reactions with massless particles and the power eight is found from dimensional consideration. Because of conservation of the sum of baryonic and leptonic numbers \(\xi_l = -\xi_q/3\).

The case of an essential variation of \(\dot{R}(t)\) is analogous to fast variation of \(\dot{\theta}(t)\) studied in our paper\(^8\). Clearly, it is much more complicated technically. Here we consider only the simple situation with quasi-stationary background and postpone more realistic time dependence of \(R(t)\) for the future work.

For small chemical potential the baryonic number density \((20)\) is equal to

$$n_B \approx \frac{g_x B_q}{6} \xi_q T^3.$$

(31)

and if the temperature adiabatically decreases in the course of the cosmological expansion, according to \(\dot{T} = -HT\), equation \((22)\) turns into

$$\dot{\xi}_q = \Gamma \left[ \frac{9f\dot{R}(t)}{10m_0^2 T} - \xi_q \right],$$

(32)
where $\Gamma \sim g^2 T^5/m_P^4$ is the rate of B-nonconserving reactions.

If $\Gamma$ is in a certain sense large, this equation can be solved in stationary point approximation as

$$\xi_q = \xi_{eq} - \frac{\dot{\xi}_{eq}}{\Gamma}, \text{ where } \xi_{eq} = \frac{9}{10} \frac{f R}{m_P^2 T}. \quad (33)$$

If we substitute $\xi_{eq}$ into Eq. (20) we arrive to the fourth order equation for $R$.

According to the comment below Eq. (20), the contribution of thermal matter into this equation can be neglected, and we arrive to the very simple fourth order differential equation:

$$\frac{d^4 R}{dt^4} = \lambda^4 R, \quad (34)$$

where $\lambda^4 = C_\lambda m_P^2 m_0^4/T^2$ with $C_\lambda = 5/(36\pi f^2 g_s B_q)$. Deriving this equation we neglected the Hubble parameter factor in comparison with time derivatives of $R$. It is justified a posteriori because the calculated $\lambda$ is much larger than $H$.

Evidently equation (34) has extremely unstable solution with instability time by far shorter than the cosmological time. This instability would lead to an explosive rise of $R$, which may possibly be terminated by the nonlinear terms proportional to the product of $H$ to lower derivatives of $R$. Correspondingly one may expect stabilization when $HR \sim \dot{R}$, i.e. $H \sim \lambda$. Since

$$\dot{H} + 2H^2 = -R/6, \quad (35)$$

$H$ would also exponentially rise together with $R$, $H \sim \exp(\lambda t)$ and $\lambda H \sim R$. Thus stabilization may take place at $R \sim \lambda^2 \sim m_P m_0^4/T$. This result should be compared with the normal General Relativity value $R_{GR} \sim T_{\text{matter}}/m_P^2$, where $T_{\text{matter}}$ is the trace of the energy-momentum tensor of matter.

4. Discussion and conclusion

For more accurate analysis numerical solution will be helpful, which we will perform in another work. The problem is complicated because the assumption of slow variation of $\dot{R}$ quickly becomes broken and the collision integral in time dependent background is not so simply tractable as the usual stationary one. The technique for treating kinetic equation in non-stationary background is presented in Ref. 8.

For evaluation of $R(t)$ in this case numerical calculations are necessary, which will be presented elsewhere. Here we describe only the basic features of the new effect of instability in gravitational baryogenesis.

To conclude we have shown that gravitational baryogenesis in the simplest versions discussed in the literature is not realistic because the instability of the emerging gravitational equations destroys the standard cosmology. Some stabilization mechanism is strongly desirable. Probably stabilization may be achieved in a version of F(R)-theory.
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References
1. A. G. Cohen and D. B. Kaplan, Phys. Lett. B 199 (1987) 251.
2. A. G. Cohen and D. B. Kaplan, Nucl. Phys. B 308 (1988) 913.
3. A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B 263 (1991) 86.
4. A. D. Dolgov, Phys. Rept. 222 (1992) 309.
5. V. A. Rubakov and M. E. Shaposhnikov, Usp. Fiz. Nauk 166 (1996) 493 [Phys. Usp. 39 (1996) 461].
6. A. Dolgov and K. Freese, Phys. Rev. D 51 (1995) 2693.
7. A. Dolgov, K. Freese, R. Rangarajan and M. Srednicki, Phys. Rev. D 56 (1997) 6155.
8. E. V. Arbuzova, A. D. Dolgov and V. A. Novikov, Phys. Rev. D 94 (2016) no.12, 123501.
9. H. Davoudiasl, R. Kitano, G. D. Kribs, H. Murayama and P. J. Steinhardt, Phys. Rev. Lett. 93 (2004) 201301.
10. E. V. Arbuzova and A. D. Dolgov, Phys. Lett. B 769 (2017) 171.
11. E. V. Arbuzova and A. D. Dolgov, JCAP 1706 (2017) no.06, 001.