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Curvature oscillations in modified gravitational baryogenesis and ultra high energy cosmic rays

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The instability problem of the gravitational baryogenesis (GBG) is analysed. It is shown that the explosive growth of the curvature scalar, inherent to GBG, can be terminated by introducing the R^2-term into the classical action of General Relativity. As a result, the exponential rising curvature is transformed into a quickly oscillating one. The high-frequency curvature oscillations lead to the production of energetic particles which, according to the estimates presented, could make a noticeable contribution to ultra high energy cosmic rays (UHECR).

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1. Introduction

Our existence critically relies on the prevalence of matter over antimatter, a fact strongly supported by observational evidence. In the local universe, matter clearly dominates. The amount of antimatter is very small, which can be explained as a result of high-energy collisions in space. The presence of large regions of antimatter nearby would lead to the production of high-energy radiation due to matter-antimatter annihilation, which is not actually observed. Any initial asymmetry during inflation could not resolve the issue of the observed excess of matter over antimatter, as the energy density associated with the baryon number did not allow inflation to last long enough.

On the other hand, matter and antimatter appear to possess similar properties, leading to the expectation of a universe symmetric in matter and antimatter. A satisfactory model of our universe should be able to explain the origin of the matter-antimatter asymmetry that surrounds us. The term "baryogenesis" is used to denote the "generation of asymmetry" between baryons (essentially protons and neutrons) and antibaryons (antiprotons and antineutrons).

In 1967, A.D. Sakharov introduced three conditions, now known as the *Sakharov principles* [1], to produce matter-antimatter asymmetry from an initially symmetric universe. These principles are: 1) non-conservation of baryonic number; 2) symmetry breaking between particles and antiparticles; 3) deviation from thermal equilibrium. However, not all of these conditions are strictly necessary. There are interesting scenarios of baryogenesis where one or several of these conditions are not satisfied. For instance, spontaneous baryogenesis (SBG) and gravitational baryogenesis (GBG) can occur in thermal equilibrium and do not require explicit C and CP violation.

The idea that cosmological baryon asymmetry can be generated by spontaneous baryogenesis in thermal equilibrium was first proposed in the original paper by A. Cohen and D. Kaplan in 1987 [2], and in subsequent papers by A. Cohen, D. Kaplan, and A. Nelson [3, 4] (for a review see [5–8]). The theory's underlying symmetry, ensuring conservation of total baryonic number in the unbroken phase, is assumed to be spontaneously broken. In the broken phase, the Lagrangian density acquires the term $\mathcal{L}_{SBG} = (\partial_\mu \theta) J^\mu_B$, which in a spatially homogeneous case simplifies to:

$$\mathcal{L}_{SBG} = \dot{\theta}\, n_B\,,\ n_B \equiv J^0_B, \tag{1}$$

Here, θ is a (pseudo) Goldstone field, and J^μ_B is the baryonic current of matter fields, becoming non-conserved. n_B is the baryonic number density, suggesting $\dot{\theta}$ could be identified with the chemical potential, μ_B, of the system. However, such identification is questionable and depends on the chosen representation for the fermionic fields [9].

Motivated by spontaneous baryogenesis, the concept of gravitational baryogenesis was proposed [10]. The SBG scenario was enhanced by introducing a coupling between the baryonic current and the derivative of the curvature scalar R:

$$S_{GBG} = -\frac{1}{M^2}\int d^4x \sqrt{-g}\,(\partial_\mu R) J^\mu_B\,, \tag{2}$$

where g is the determinant of the space-time metric tensor, and the mass parameter M defines the energy scale of baryogenesis. There are many articles on this subject, with a partial list of references included in [11–15]. According to these papers, the GBG mechanism can successfully explain the magnitude of the cosmological baryon asymmetry of the universe.

However, as shown in [16, 17], the addition of a curvature-dependent term to the standard Lagrangian of general relativity (GR) leads to higher-order gravitational equations of motion, which become highly unstable in response to small perturbations. In the presented paper, we demonstrate that the stability problem can be resolved by introducing a curvature-squared term into the Hilbert-Einstein action of GR.

2. Instability problem of gravitational baryogenesis

2.1 Bosonic case

Let's examine the model in which the baryonic number is conveyed by a complex scalar field ϕ with potential $U(\phi, \phi^*)$:

$$S = -\int d^4x\,\sqrt{-g}\left[\frac{M_{Pl}^2}{16\pi}R + \frac{1}{M^2}(\partial_\mu R)J^\mu - g^{\mu\nu}\partial_\mu\phi\,\partial_\nu\phi^* + U(\phi,\phi^*)\right] + S_m\,, \tag{3}$$

where $M_{Pl} = 1.22\cdot 10^{19}$ GeV is the Planck mass, S_m is the action for matter.

If the potential $U(\phi)$ is not invariant with respect to the $U(1)$-rotation, $\phi \to e^{i\beta}\phi$, the baryonic current defined in the usual way

$$J_\mu = iq(\phi^*\partial_\mu\phi - \phi\partial_\mu\phi^*)$$

is not conserved. Here q represents the baryonic number of field ϕ.

The equation of motion for the curvature scalar in this model is expressed as follows:

$$\frac{M_{Pl}^2}{16\pi}R + \frac{1}{M^2}\left[(R+3D^2)D_\alpha J^\alpha + J^\alpha\,D_\alpha R\right] - D_\alpha\phi\,D^\alpha\phi^* + 2U(\phi) = -\frac{1}{2}T^\mu_\mu\,, \tag{4}$$

where D_μ is the covariant derivative, T^μ_μ is the trace of the energy-momentum tensor of matter derived from the action S_m.

For a homogeneous curvature scalar $R(t)$ in a spatially flat Friedmann-Lemaitre-Robertson-Walker (FLRW) metric, $ds^2 = dt^2 - a^2(t)d\mathbf{r}^2$, Eq. (4) simplifies to:

$$\frac{M_{Pl}^2}{16\pi}R + \frac{1}{M^2}\left[(R + 3\partial_t^2 + 9H\partial_t)D_\alpha J^\alpha + \dot{R}\,J^0\right] = -\frac{T^{(tot)}}{2}. \tag{5}$$

Here J^0 is the baryonic number density of the ϕ-field, $H = \dot{a}/a$ is the Hubble parameter, $T^{(tot)}$ is the trace of the energy-momentum tensor of matter including contribution from the ϕ-field.

In the homogeneous case, the covariant divergence of the current is given by:

$$D_\alpha J^\alpha = \frac{2q^2}{M^2}\left[\dot{R}\,(\phi^*\dot{\phi} + \phi\dot{\phi}^*) + (\ddot{R} + 3H\dot{R})\,\phi^*\phi\right] + iq\left(\phi\frac{\partial U}{\partial\phi} - \phi^*\frac{\partial U}{\partial\phi^*}\right). \tag{6}$$

The expectation values of the products of quantum operators ϕ, ϕ^* and their derivatives after the thermal averaging, according to [16], are equal to:

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

where T is the plasma temperature.

The equation of motion for the classical field R in the cosmological plasma is as follows:

$$\frac{M_{Pl}^2}{16\pi} R + \frac{q^2}{6M^4} \left(R + 3\partial_t^2 + 9H\partial_t \right) \left[\left(\ddot{R} + 3H\dot{R} \right) T^2 \right] + \frac{1}{M^2} \dot{R} \left\langle J^0 \right\rangle = -\frac{T^{(tot)}}{2} . \tag{7}$$

Here $\langle J^0 \rangle$ is the thermally averaged value of the baryonic number density of ϕ. This term can be neglected, as it is initially small and later becomes subdominant.

Focusing only on the linear in curvature terms and disregarding higher powers of R, such as R^2 or HR, we derive the fourth-order differential equation for the curvature scalar:

$$\frac{d^4 R}{dt^4} + \mu^4 R = -\frac{1}{2} T^{(tot)}, \quad \mu^4 = \frac{M_{Pl}^2 M^4}{8\pi q^2 T^2} . \tag{8}$$

The homogeneous part of this equation has exponential solutions:

$$R \sim e^{\lambda t}, \quad \lambda = |\mu| e^{i\pi/4 + i\pi n/2}, \quad n = 0, 1, 2, 3. \tag{9}$$

There are two solutions with positive real parts of λ. This implies that the curvature scalar is exponentially unstable with respect to small perturbations. Consequently, R should rise exponentially fast over time and oscillate rapidly around this ascending function.

2.2 Fermionic case

We now turn to a more realistic scenario where the baryonic number is carried by fermions. The action for this case is presented as follows:

$$S = -\int d^4 x \sqrt{-g} \left[\frac{M_{Pl}^2}{16\pi} R - \mathcal{L}[Q, L] \right] + S_m, \tag{10}$$

where the Lagrangian $\mathcal{L}[Q, L]$ is given by

$$\begin{aligned} \mathcal{L}[Q, L] = \frac{i}{2} (\bar{Q} \gamma^\mu \nabla_\mu Q - \nabla_\mu \bar{Q}\, \gamma^\mu Q) - m_Q \bar{Q}\, Q + \frac{i}{2} (\bar{L} \gamma^\mu \nabla_\mu L - \nabla_\mu \bar{L} \gamma^\mu L) - m_L \bar{L}\, L \\ + \frac{g}{m_X^2} \left[(\bar{Q}\, Q^c)(\bar{Q} L) + (\bar{Q}^c Q)(\bar{L} Q) \right] + \frac{d}{M^2} (\partial_\mu R) J^\mu_Q . \end{aligned} \tag{11}$$

In this setup: Q represents the quark (or quark-like) field with a non-zero baryonic number B_Q, L is another fermionic field, possibly a lepton. ∇_μ denotes the covariant derivative of Dirac fermions in the tetrad formalism. $J^\mu_Q = B_Q \bar{Q} \gamma^\mu Q$ is the quark current, where γ^μ are the curved space gamma-matrices; $d = \pm 1$ is a dimensionless coupling constant introduced to allow for an arbitrary sign. The four-fermion interaction between quarks and leptons is included to ensure the necessary non-conservation of the baryonic number.

Taking trace of gravitational equations for matter, we come to the following:

$$\begin{aligned} -\frac{M_{Pl}^2}{8\pi} R = m_Q \bar{Q} Q + m_L \bar{L} L + \frac{2g}{m_X^2} \left[(\bar{Q}\, Q^c)(\bar{Q} L) + (\bar{Q}^c Q)(\bar{L} Q) \right] - \frac{2d}{M^2} (R + 3D^2) D_\alpha J^\alpha_Q \\ + T_{other} , \end{aligned} \tag{12}$$

where T_{other} is the trace of the energy-momentum tensor of all other fields. At the relativistic stage, we can put $T_{other} = 0$.

Using the kinetic equation, we find an explicit dependence of $D_\alpha J^\alpha_Q$ on $\dot{R}$, in cases where the current is not conserved, as detailed in Ref. [17]. Consequently, this results in a high-order (fourth-order) equation for R.

When disregarding the contribution from thermal matter and considering the FLRW-background, the following equation emerges:

$$\frac{M_{Pl}^2}{8\pi} R = \frac{2d}{M^2}(R + 3D^2)(\partial_t + 3H) n_B \,. \tag{13}$$

The baryonic number density is derived from the kinetic equation:

$$n_B \sim \frac{9d}{10} \frac{g_s B_Q \dot{R}}{M^2 T}, \tag{14}$$

where g_s denotes the number of quark spin states.

Neglecting the H-factor in comparison to the time derivatives of R, we are led to a very simple fourth-order differential equation for the curvature scalar:

$$\frac{d^4 R}{dt^4} + \lambda^4 R = 0, \quad \lambda^4 = \frac{5 M_{Pl}^2 M^4}{36\pi g_s B_Q^2 T^2}. \tag{15}$$

This equation yields an extremely unstable solution, with the instability time being significantly shorter than the cosmological time. Such a situation would result in an explosive increase of R.

3. Stabilization of GBG in R^2-modified gravity

As shown in the previous section, traditional gravitational baryogenesis, which is based on the assumption of an interaction between the derivative of the curvature scalar, $\partial_\mu R$, and the baryonic current, J^μ_B, can successfully explain the magnitude of the cosmological baryon asymmetry of the universe. However, the back-reaction of the created nonzero baryonic density leads to strong instability in the cosmological evolution.

A potential stabilization mechanism could be realized in R^2- modified gravity, as described by the following action:

$$S_{Grav} = -\frac{M_{Pl}^2}{16\pi} \int d^4x \sqrt{-g} \left(R - \frac{R^2}{6M_R^2} \right). \tag{16}$$

The inclusion of the R^2-term in the standard action of General Relativity arises from one-loop corrections to the energy-momentum tensor of matter in curved space-time, initially identified in Ref. [18]. This approach was further developed by V. T. Gurovich and A. A. Starobinsky [19]. It is important to note that the R^2-term leads to the excitation of a scalar degree of freedom, known as *scalaron*, where M_R represents the scalaron's mass. In the very early universe, the R^2-term can induce inflation [20] and affect density perturbations. The observed amplitude of these density perturbations suggests that $M_R = 3 \cdot 10^{13}$ GeV [21].

Below we consider the models from sections 2.1 and 2.2 and show how the instability problem of GBG can be solved in R^2-modified gravity [22].

3.1 Stabilization: bosonic case

We consider action (3) with the addition of an R^2-term introduced for stabilization purposes:

$$S_{tot}[\phi] = -\int d^4x\,\sqrt{-g}\left[\frac{M_{Pl}^2}{16\pi}\left(R - \frac{R^2}{6M_R^2}\right) + \frac{1}{M^2}(\partial_\mu R)J^\mu_{(\phi)} - g^{\mu\nu}\partial_\mu\phi\,\partial_\nu\phi^* + U(\phi,\phi^*)\right] + S_m. \quad (17)$$

Equation (4) for the curvature evolution correspondingly takes the form:

$$\frac{M_{Pl}^2}{16\pi}\left(R + \frac{1}{M_R^2}D^2R\right) + \frac{1}{M^2}\left[(R+3D^2)D_\alpha J^\alpha_{(\phi)} + J^\alpha_{(\phi)}\,D_\alpha R\right] - D_\alpha\phi\,D^\alpha\phi^* + 2U(\phi)$$
$$= -\frac{1}{2}T^{(matt)}. \quad (18)$$

For relativistic matter $T^{(matt)} = 0$.

In a spatially flat FLRW-metric the equation is as follows:

$$\frac{M_{Pl}^2}{16\pi}\left[R + \frac{1}{M_R^2}(\partial_t^2 + 3H\partial_t)R\right] + \frac{1}{M^2}\left[(R + 3\partial_t^2 + 9H\partial_t)D_\alpha J^\alpha_{(\phi)} + \dot R\,J^0_{(\phi)}\right]$$
$$+2U(\phi) - (D_\alpha\phi)(D^\alpha\phi^*) = 0. \quad (19)$$

Upon substituting the divergence of the current, $D_\alpha J^\alpha_{(\phi)}$, using Eqs. (6) and (7), we derive the 4th-order differential equation for the evolution of the curvature scalar:

$$\frac{M_{Pl}^2}{16\pi}\left(R + \frac{1}{M_R^2}D^2R\right) + \frac{q^2}{6M^4}\left(R + 3\partial_t^2 + 9H\partial_t\right)\left[\left(\ddot R + 3H\dot R\right)T^2\right] + \frac{1}{M^2}\dot R\,\langle J^0_{(\phi)}\rangle = -\frac{T^\mu_\mu(\phi)}{2}. \quad (20)$$

Focusing only on the dominant terms, we can simplify the preceding equation to:

$$\frac{d^4R}{dt^4} + \frac{\kappa^4}{M_R^2}\frac{d^2R}{dt^2} + \kappa^4 R = -\frac{1}{2}T^\mu_\mu(\phi), \quad \kappa^4 = \frac{M_{Pl}^2M^4}{8\pi q^2T^2}. \quad (21)$$

This equation is similar to Eq. (8), yet it is distinguished by a stabilizing term proportional to the second derivative of curvature.

3.2 Stabilization: fermionic case

In a manner completely analogous to section 3.1, we consider the stabilization of GBG with fermions by incorporating an R^2-term into action (10). The action now is represented as:

$$S_{tot}[Q,L] = -\int d^4x\,\sqrt{-g}\left[\frac{M_{Pl}^2}{16\pi}\left(R - \frac{R^2}{6M_R^2}\right) + \frac{d}{M^2}(\partial_\mu R)J^\mu_Q - \mathcal{L}[Q,L]\right], \quad (22)$$

where $\mathcal{L}[Q,L]$ is given by Eq. (12).

The equation for the evolution of curvature takes the form:

$$-\frac{M_{Pl}^2}{8\pi}\left(R + \frac{1}{M_R^2}D^2R\right) = \quad (23)$$
$$m_Q\bar QQ + m_L\bar LL + +\frac{2g}{m_X^2}\left[(\bar Q\,Q^c)(\bar QL) + (\bar Q^cQ)(\bar LQ)\right] - \frac{2d}{M^2}(R+3D^2)D_\alpha J^\alpha_Q + T_{matt}\,.$$

In the early universe, when various species are relativistic, $T_{matt} = 0$. A higher order equation for R emerges after substituting the current divergence $D_\alpha J^\alpha_Q$, which is calculated from the kinetic equation in the external field R.

In complete analogy with the previous cases, we derive the following equation:

$$\frac{d^4R}{dt^4} + \frac{\kappa_f^4}{M_R^2}\frac{d^2R}{dt^2} + \kappa_f^4 R = 0, \quad \kappa_f^4 = \frac{5M_{Pl}^2 M^4}{36\pi g_s B_Q^2 T^2}. \tag{24}$$

The value of κ_f differs only slightly from κ in the scalar case (21).

3.3 Stability condition

Searching for a solution to Eq. (21) in the form $R = R_{in}\exp(\lambda t)$, we arrive at the characteristic equation:

$$\lambda^4 + \frac{\kappa^4}{M_R^2}\lambda^2 + \kappa^4 = 0 \tag{25}$$

with the eigenvalues λ determined by the expression:

$$\lambda^2 = -\frac{\kappa^4}{2M_R^2} \pm \kappa^2\sqrt{\frac{\kappa^4}{4M_R^4} - 1}. \tag{26}$$

Instability is absent if $\lambda^2 < 0$, and Eq. (21) then has only oscillating solutions. This is realized when $\kappa^4 > 4M_R^4$. Using the expression for κ^4 from Eq. (21) and taking $M_R = 3\cdot 10^{13}$ GeV, we establish the stability condition:

$$M > 3\cdot 10^4\,\text{GeV}\left(\frac{qT}{\text{GeV}}\right)^{1/2}, \tag{27}$$

which is valid for all significant values of M.

The value of λ depends on the relation between κ and M_R. If $\kappa \sim M_R$ then the frequency of curvature oscillations is of the order of M_R and $|\lambda| \sim M_R$. If $\kappa \gg M_R$ then there are two possible solutions: $|\lambda| \sim M_R$ and $|\lambda| \sim \kappa(\kappa/M_R) \gg M_R$.

Since the value of κ_f in Eq. (24) only slightly differs numerically from κ in Eq. (21) and has the same dependence on the essential parameters, the solutions of Eqs. (21) and (24) are practically identical. Thus, the stability condition (27) is applicable to both the bosonic and fermionic cases.

High frequency oscillations of R would lead to effective gravitational particle production and subsequently to the damping of these oscillations. It is intriguing to consider how such high-frequency curvature oscillations might contribute to the ultra-high-energy cosmic ray (UHECR) spectrum.

4. Curvature oscillations and high energy cosmic rays

We posit that superheavy dark matter (DM) particles were generated by the oscillating curvature scalar $R(t)$ within the framework of Starobinsky inflation [20], as described by action (16). This concept has been explored in detail in the works [23–25]. In the R^2-theory the oscillating curvature can be interpreted as an effective scalar field, scalaron, with mass M_R and decay width Γ. We focus on the scenario where scalaron decays yield particles with masses around 10^{21} eV, corresponding to the energy range of UHECR. Let us stress, that the superheavy particle decays are suggested as source only for the CR with energies above 10^{20} eV, that cannot be explained by the canonical astrophysical processes. Such decays have been suggested in multitude of papers. For a review see Ref. [26]. There are only several events with $E > 10^{20}$ eV and statistics is too low to be inconsistent with the data. The contribution of the suggested mechanism into lower energy cosmic rays is sufficiently small, so it doesn't distort the obsereved flux.

Dark matter particles are typically assumed to be absolutely stable. Nevertheless, Ya. B. Zeldovich proposed a mechanism [27, 28], that implies the decay of any presumably stable particle through the creation and evaporation of virtual black holes.

In our work [29], we demonstrate that superheavy DM particles with masses about 10^{12} GeV may decay through the virtual black hole, with a life-time that is only a few orders of magnitude longer than the age of the universe. The decay of such particles could significantly contribute to the UHECR spectrum. This scenario becomes plausible in theories where gravitational coupling increases at small distances or high energies.

We examine the model suggested in Refs. [30, 31], where the observable universe, containing Standard Model particles, is confined to a 4-dimensional brane within a (4+d)-dimensional bulk. In this framework, gravity is not restricted to the brane but propagates throughout the entire bulk. As a result, the Planck mass M_{Pl} in such scenarios is reinterpreted as an effective long-distance 4-dimensional parameter. The relationship between M_{Pl} and the fundamental gravity scale M_* is established as follows:

$$M_{Pl}^2 \sim M_*^{2+d} R_*^d, \tag{28}$$

where R_* is the size of the extra dimensions:

$$R_* \sim \frac{1}{M_*}\left(\frac{M_{Pl}}{M_*}\right)^{2/d}. \tag{29}$$

For future application, we have chosen $M_* \approx 3 \times 10^{17}$ GeV, resulting in $R_* \sim 10^{(4/d)}/M_* > 1/M_*$.

Analogous to the proton decay, $p \to l^+ \bar{q} q$, let us consider the folliwng decay of X-particle: $X \to L^+ \bar{q}_* q_*$, as depicted in the diagram in Fig. 1.

According to the calculations in Ref. [32], the decay width of the proton into a positively charged lepton and a quark-antiquark pair is given by:

$$\Gamma(p \to l^+ \bar{q} q) = \frac{m_p\, \alpha^2}{2^{12}\, \pi^{13}} \left(\ln \frac{M_{Pl}^2}{m_q^2}\right)^2 \left(\frac{\Lambda}{M_{Pl}}\right)^6 \left(\frac{m_p}{M_{Pl}}\right)^{4+\frac{10}{d+1}} \int_0^{1/2} dx x^2 (1-2x)^{1+\frac{5}{d+1}}, \tag{30}$$

where $m_p \approx 1$GeV is the proton mass, $m_q \sim 300$ MeV is the constituent quark mass, $\Lambda \sim 300$ MeV is the QCD scale parameter, $\alpha = 1/137$ is the fine structure constant, and d is the number of "small' extra dimensions. The QCD coupling constant α_s is supposed to be equal to unity. It can be verified that the proton decay rate is exceedingly small. This results in a corresponding life-time of 7.3×10^{198} years which is significantly longer than the age of the universe, $t_U \approx 1.5 \times 10^{10}$ years.

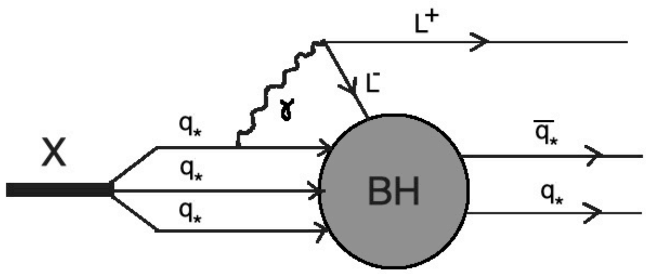

Figure 1: Diagram describing X-particle decay into $L^+\bar{q}_*q_*$ through virtual black hole.

The mentioned case of a decaying proton serves merely as an illustrative example. Our primary focus is on superheavy dark matter (DM) particles with masses around 10^{12} GeV. We aim to develop a scenario where these superheavy DM particles, decaying through a virtual black hole, have a life-time only a few orders of magnitude longer than the age of the universe.

We consider the decay process $X \to L^+\bar{q}_*q_*$ and assume that heavy dark matter X-particle, with a mass $M_X \sim 10^{12}$ GeV, is composed of three heavy quarks, q_*, each having a comparable mass. In this scenario, we leave Λ_* as a free parameter. The life-time of X-particles can be estimated using Eq. (30) with modifications to account for the characteristics of the X-particle. Specifically, we replace the fine structure constant $\alpha = 1/137$ with $\alpha_* = 1/50$, the proton mass m_p with $M_X = 10^{12}$ GeV, and the constituent quark mass m_{q_*} with 10^{12} GeV. Additionally, we consider $d = 7$ extra dimensions in this calculation. Thus, we obtain:

$$\tau_X = \frac{1}{\Gamma_X} \approx 6.6 \times 10^{-25}\mathrm{s} \cdot \frac{2^{10}\pi^{13}}{\alpha_*^2} \left(\frac{\mathrm{GeV}}{M_X}\right) \left(\frac{M_*}{\Lambda_*}\right)^6 \left(\frac{M_*}{M_X}\right)^{4+\frac{10}{d+1}} \left(\ln \frac{M_*}{m_{q_*}}\right)^{-2} I(d)^{-1}, \tag{31}$$

where we took 1/GeV = 6.6×10^{-25}s and

$$I(d) = \int_0^{1/2} dx x^2 (1-2x)^{1+\frac{5}{d+1}}, \quad I(7) \approx 0.0057. \tag{32}$$

Now all the parameters, except for Λ_*, are fixed: $M_* = 3 \times 10^{17}$ GeV, $M_X = 10^{12}$ GeV, $m_{q_*} \sim M_X$, allowing us to estimate the life-time of X-particles as follows:

$$\tau_X \approx 7 \times 10^{12} \text{ years} \left(10^{15}\,\mathrm{GeV}/\Lambda_*\right)^6 \quad vs \quad t_U \approx 1.5 \times 10^{10} \text{ years}. \tag{33}$$

By slightly adjusting Λ_* to near 10^{15} GeV, we can set the life-time of dark matter X-particles within an interesting range. Such particles would be sufficiently stable to act as cosmological dark matter, and their decays could significantly contribute to cosmic rays at ultra-high energies. Anyway, the lifetime is a free parameter that can be adjusted to satisfy the existing constraints, which are valid in the lower energy range.

5. Conclusions

In the gravitational baryogenesis scenario, the derivative coupling of the baryonic current to the curvature scalar results in fourth-order equations for the gravitational field. These equations exhibit instability with respect to small perturbations of the FLRW-background, leading to an exponential increase in curvature. For a wide range of cosmological temperatures, the development of instability occurs much faster than the rate of the Universe's expansion. The stability issue can be addressed by incorporating an R^2-term into the Hilbert-Einstein action, which induces oscillations of the curvature and leads to efficient particle production.

In the model of modified high-dimensional gravity, there may exist superheavy dark matter (DM) particles that are stable against conventional particle interactions. Nonetheless, these DM particles are expected to decay through the formation of virtual black holes. With an appropriate selection of parameters, the lifetime of such quasi-stable particles may exceed the age of the universe by only 3-4 orders of magnitude. This allows X-particles to significantly contribute to the flux of ultra high-energy cosmic rays with energies above 10^{20} eV, not distorting the lower energy bulk. The suggested mechanism could enable the efficient generation of high-energy cosmic ray neutrinos, as detected by observatories such as IceCube and Baikal.

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