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Standard model baryon number violation seeded by black holes

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ABSTRACT

We show that black holes with a Schwarzschild radius of the order of the electroweak scale may act as seeds for the baryon number violation within the Standard model via sphaleron transitions. The corresponding rate is faster than the one in the pure vacuum and baryon number violation around black holes can take place during the evolution of the universe after the electroweak phase transition. We show however that this does not pose any threat for a pre-existing baryon asymmetry in the universe.

1. Introduction

It is well-known that, within the Standard Model (SM) of electroweak interactions, the baryon (B) and the lepton (L) symmetries are accidental and it is not possible to violate their corresponding charges at any order of perturbation theory. Nevertheless, non-perturbative effects may give rise to processes which violate the baryon and the lepton numbers. Indeed, the presence of the non-abelian group $SU(2)_L$ within the SM gauge group implies that the ground state is the sum of an infinite number of vacua which are classically degenerate and have different baryon (and lepton) numbers. Static configurations, called sphalerons [1], corresponding to unstable solutions of the equations of motion and to saddle points of the energy functional, interpolate between two nearby vacua.

The probability of baryon number violation to occur in the vacuum through sphaleron transitions is exponentially suppressed [2]

$$\Gamma_B \sim e^{-4\pi/\alpha_{\rm W}} \sim e^{-150}, \tag{1}$$

where $\alpha_{\rm W} = g_2^2/4\pi$ is the $SU(2)_L$ gauge coupling constant. Such an exponential factor is interpreted as the probability of making a transition from one classical vacuum to the closest one by quantum tunneling, going through a barrier of energy $E_{\rm sph} \sim 10$ TeV thanks to the formation of a sphaleron. In more extreme situations, like the primordial Universe, baryon and lepton number violation processes may be however faster through classical transitions induced by the high-temperature environment and play a significant role in the generation of the baryon asymmetry [3].

There are also arguments suggesting that all global symmetries, including the baryon one, are violated when including gravity [4]. In particular, no-hair theorems tell us that global charges are swallowed by Black Holes (BHs). Indeed, quanta with global charge may scatter with a BH, leaving behind a BH with a slightly larger mass, but indeterminate global charge as dictated by the no-hair theorem. At the level of effective field theory, one can imagine to integrate out virtual BH states of mass $M_{\rm BH}$ arising from quantum gravity, leading to higher-dimensional baryon number violating operators suppressed by powers of $M_{\rm BH}$, where $M_{\rm BH}$ might be as small as the Planck mass $M_{\rm Pl}$.

What about baryon number violation induced by sphaleron transitions in the presence of BHs? In general, tunneling processes may be catalysed by the presence of impurities. A BH is a gravitational impurity and indeed it has been shown that BHs can trigger electroweak SM vacuum instability in their vicinity, both at zero temperature [5–8] and in the early universe [9–14], and baryon number violations through interactions with skyrmions [15,16].

Since we are dealing with SM sphaleron configurations, a simple estimate tells us that the typical Schwarzschild radius of the BH able to alter the rate of the baryon number violation is

$$r_{\rm S} = 2GM_{\rm BH} \sim \frac{1}{M_{\rm W}} \sim \frac{1}{g_2 v}, \tag{2}$$

where $G = 1/M_{\rm Pl}^2$ and $v = 246$ GeV is the Vacuum Expectation Value (VEV) of the Higgs field. This leads to BH masses in the ballpark of

$$M_{\rm BH} = \mathcal{O}(1) \cdot 10^{-22} M_\odot \sim 10^{17} M_{\rm Pl}, \tag{3}$$

i.e. to BHs which evaporate with a typical lifetime of $\mathcal{O}(1)$ yr and which might have been present during the evolution of the Universe.

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We are going to show that baryon number violation through sphaleron transitions in the presence of such BHs can be faster than in the pure vacuum and we will offer as well some considerations about what may happen should these tiny BHs be present during the evolution of the universe.

2. Baryon number violation seeded by BHs

To study the influence of BHs on the sphaleron transitions we start from the action of the Higgs doublet field ϕ along with a $SU(2)_L$ gauge field W^a_μ (including the abelian hypercharge group $U(1)_Y$ does not change our results) in a curved spacetime

$$S=\int_{\mathcal{M}} \mathrm{d}^4x\sqrt{-g}\left[\frac{\mathcal{R}}{16\pi G}-g^{\mu\nu}(D_\mu\phi)^\dagger D_\nu\phi-V(\phi)\right.$$
$$\left.-\frac{1}{4}g^{\mu\rho}g^{\nu\sigma}F^a_{\rho\sigma}F^a_{\mu\nu}\right]+\frac{1}{8\pi G}\int_{\partial\mathcal{M}}\mathcal{K}\sqrt{\gamma}\mathrm{d}^3y, \tag{4}$$

where $V(\phi)$ is the Higgs potential and we have added the Gibbons-Hawking-York boundary term as we deal with a spacetime manifold $\mathcal{M}$ with a BH horizon. The spacetime geometry around the BH can be taken static and spherically symmetric, such that its metric takes a Schwarzschild-like form

$$\mathrm{d}s^2=-e^{2\delta(r)}A(r)\mathrm{d}t^2+A^{-1}(r)\mathrm{d}r^2+r^2\mathrm{d}\Omega^2,$$
$$A(r)=1-\frac{2GM(r)}{r}, \tag{5}$$

where $A(r)$ vanishes at the horizon

$$r_S\equiv 2GM(r_S)=2GM_{\rm BH}. \tag{6}$$

A suitable ansatz for the gauge and Higgs field is [1]

$$W^a_i\sigma^a\mathrm{d}x^i=-\frac{2i}{g_2}f(g_2vr)\mathrm{d}U^\infty(U^\infty)^{-1},$$
$$\phi=\frac{v}{\sqrt{2}}h(g_2vr)U^\infty\begin{pmatrix}0\\1\end{pmatrix}, \tag{7}$$

where

$$U^\infty=\frac{1}{r}\begin{pmatrix}z & x+iy\\ -x+iy & z\end{pmatrix}. \tag{8}$$

Since we are ultimately interested in the energy functional, we perform an analytical continuation of the action to the Euclidean metric with $t=i\tau$, taking τ to be periodic with period $1/T$ (to be identified with the relevant temperature of the system).

By setting $\xi=g_2vr$ and expanding the mass with respect to its value at the horizon

$$M(\xi)=M_{\rm BH}+\delta M(\xi), \tag{9}$$

we can write the equations of motion

$$\delta M'=\frac{4\pi v}{g_2}\left[\frac{\xi^2}{2}A(h')^2+\epsilon\frac{\xi^2}{4}(h^2-1)^2+4A(f')^2\right.$$
$$\left.+\frac{8}{\xi^2}f^2(1-f)^2+\xi^2\mathcal{F}(\xi)\tilde\epsilon h^2+h^2(1-f)^2\right],$$
$$\delta'=2\alpha^2\left[\frac{8}{\xi}(f')^2+\xi(h')^2\right],$$
$$(Ae^\delta f')'=\frac{2}{\xi^2}e^\delta f(1-f)(1-2f)-\frac{1}{4}e^\delta h^2(1-f),$$
$$(\xi^2Ae^\delta h')'=\epsilon\xi^2e^\delta h(h^2-1)+2e^\delta h(1-f)^2$$
$$+2e^\delta\xi^2\mathcal{F}(\xi)\tilde\epsilon h, \tag{10}$$

where $\alpha=\sqrt{4\pi G}(v/\sqrt{2})$, $\epsilon=\lambda/g_2^2$ and $\tilde\epsilon=\tilde\lambda/768\pi^2$. Here λ is the quartic coupling of the Higgs and we have written the Higgs potential as

$$V(h,\xi)\simeq\frac{\lambda}{4}v^4(h^2-1)^2+\frac{\tilde\lambda g_2^2}{768\pi^2}\mathcal{F}(\xi)v^4h^2. \tag{11}$$

The second term is due to the vacuum polarization effect of the Hawking radiation originating at one-loop from the interactions of the Higgs with the other SM particles in the vicinity of the horizon of the BH [17]. This term is very similar to the finite temperature correction to the mass squared of the Higgs $\sim T^2h^2$ in a plasma at finite temperature T. The key difference is that the effective temperature depends on the distance from the horizon [18,19] (being $\xi_S\equiv g_2vr_S$ the dimensionless BH horizon)

$$\mathcal{F}(\xi)=\begin{cases}\frac{3}{4\xi_S^2} & \xi\simeq\xi_S,\\ \frac{1}{\xi^2} & \xi\gg\xi_S,\end{cases} \tag{12}$$

so that, close to the horizon, the correction to the potential acquires the familiar form $T_H^2h^2$, where [20]

$$T_H=\frac{1}{8\pi GM_{\rm BH}}\simeq 10\left(\frac{M_{\rm BH}}{5\cdot 10^{-22}\,M_\odot}\right)^{-1}\,\mathrm{GeV} \tag{13}$$

is the Hawking temperature. We adopt here the Unruh vacuum [21] as the most appropriate vacuum for our physical situation. Indeed, in the following we will consider the case in which the temperature of the universe is different from the Hawking temperature. As such, the Hartle-Hawking vacuum [22] is not the proper one as it assumes full and static thermal equilibrium with the surrounding plasma.

The effective coupling $\tilde\lambda$ is given by

$$\tilde\lambda=24\left(\frac{3}{16}g_2^2+\frac{1}{16}g_1^2+\frac{1}{4}y_t^2+\frac{\lambda}{2}\right)\sim 9.56, \tag{14}$$

computed in terms of the g_2, g_1 (the gauge coupling of the $U(1)_Y$ group), and top Yukawa coupling y_t, all evaluated at the electroweak scale [23].

Since $\alpha\ll 1$, one can approximate $\delta'\simeq 0$ and, given that the metric has to approach the Minkowski spacetime at infinity, the leading order solution of Eq. (10) gives $\delta\simeq 0$. The equations for the gauge and Higgs fields then simplify to

$$f''+\frac{\xi_S}{\xi(\xi-\xi_S)}f'=\frac{2}{\xi(\xi-\xi_S)}f(1-f)(1-2f)$$
$$-\frac{1}{4}\frac{\xi}{\xi-\xi_S}h^2(1-f),$$
$$h''+\left(\frac{2}{\xi}+\frac{\xi_S}{\xi(\xi-\xi_S)}\right)h'=\frac{\xi}{\xi-\xi_S}\epsilon h(h^2-1)+\frac{2\xi\mathcal{F}(\xi)}{(\xi-\xi_h)}\tilde\epsilon h$$
$$+\frac{2}{\xi(\xi-\xi_S)}h(1-f)^2. \tag{15}$$

In order to solve the equations of motion, we have to impose proper boundary conditions. At infinity the metric has to approach the Minkowski spacetime and the fields have to be in their true vacuum,

$$f(\xi)\to 1,\ \ h(\xi)\to 1,\ \ \text{when}\ \ \xi\to\infty. \tag{16}$$

At the BH horizon $\xi\to\xi_S$, one can impose the boundary conditions setting the fields in the false vacuum

$$f(\xi)\to 0,\ \ h(\xi)\to 0,\ \ \text{when}\ \ \xi\to\xi_S. \tag{17}$$

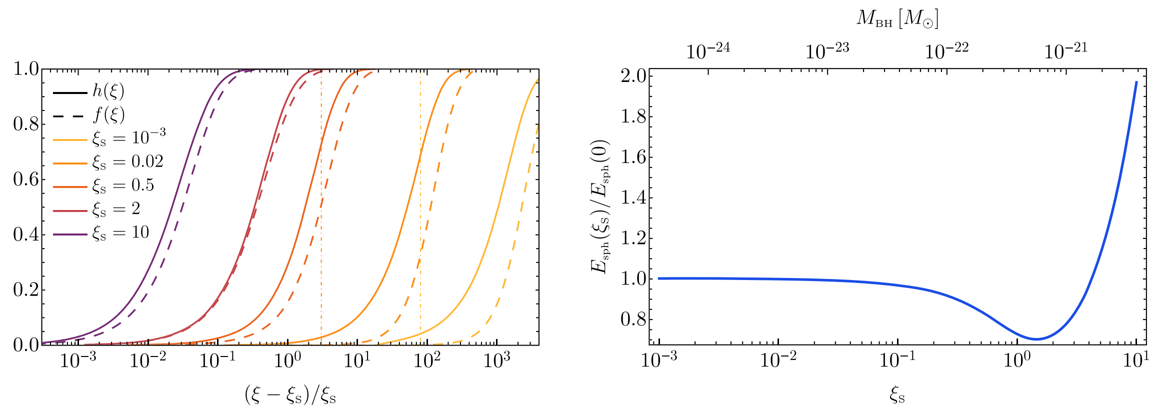

Fig. 1. *Left*: Behaviour of the Higgs and $SU(2)_L$ fields in terms of the radial coordinate for different values of the rescaled horizon ξ_S. The vertical lines indicate the radial position where the Hawking temperature is able to restore the symmetry. *Right*: Values of the rescaled sphaleron energy for different BH masses.

The numerical solutions of the equations of motion can be found in the left panel of Fig. 1 for different rescaled BH horizons. For small enough BHs, there exists a critical radius below which the vacuum polarization effect induced by the Hawking radiation leads to the restoration of the symmetry close to the horizon, nevertheless allowing for a sphaleron solution interpolating between the unbroken and broken phase.

The characteristic mass contribution at infinity is

$$\delta M_\infty = \frac{4\pi v}{g_2}\int_{\xi_S}^{\infty} \mathrm{d}\xi \left[\frac{1}{2}\xi(\xi-\xi_S)(h')^2 + \epsilon\frac{\xi^2}{4}(h^2-1)^2 + \frac{4}{\xi}(\xi-\xi_S)(f')^2 + \frac{8}{\xi^2}f^2(1-f)^2 + \xi^2\mathcal{F}(\xi)\tilde{\epsilon}h^2 + h^2(1-f)^2\right], \tag{18}$$

which has to be thought as the sphaleron energy in the presence of a BH,

$$\delta M_\infty = E_{\rm sph}(M_{\rm BH}). \tag{19}$$

Indeed, in the limit of flat spacetime with no BH ($r_S \ll 1/g_2 v$), we get

$$E_{\rm sph}(0) \simeq 1.92\frac{4\pi v}{g_2}, \tag{20}$$

which reproduces the standard result for the current physical mass of the Higgs (i.e. for $\epsilon \simeq 0.3$), see right panel of Fig. 1. Notice that the effect of the vacuum polarization in the Higgs potential, in the limit of tiny BH masses, is minor because the radius of the sphaleron configuration is located away from the Schwarzschild radius.

For small BH masses the sphaleron radius is large compared to the Schwarzschild radius and its energy is only slightly perturbed compared to the vacuum solution. As the seed BH masses increase, the sphaleron radius approaches the horizon and the BH helps catalysing the sphaleron transitions. For larger BH masses, it is energetically more costly to generate the sphaleron solution as its characteristic size is required to be larger than the BH horizon and therefore larger than $\sim 1/g_2 v$. Notice also that the minimum BH mass is consistent with the estimate (3).

3. Rate of baryon number violation seeded by BHs

How fast can the baryon number violation take place in the vicinity of a BH? The vacuum decay rate takes the form [24]

$$\Gamma_{\rm sph}(M_{\rm BH}) \sim \sqrt{\frac{B}{2\pi}}\frac{1}{\ell_{\rm sph}}e^{-B}, \tag{21}$$

where $\ell_{\rm sph}$ is the size of the sphaleron configuration. For large BH masses, it turns out to be comparable to the Schwarzschild radius r_S, while for small BH masses it is of the order of $1/g_2 v$. The dimensionless term $\sqrt{B/2\pi}$ comes from the normalization of the zero mode associated with time translation symmetry, in terms of the exponent B given by the difference between the Euclidean action of the bounce solution and the one before the transition. For a static solution this coincides with the difference between the BH areas at the horizon and at infinity. As one can easily check, for the static solution the bulk part of the energy functional vanishes due to the Hamiltonian constraint, while the boundary terms give the BH Bekenstein-Hawking entropy at the BH horizon [24]

$$B = -\frac{\mathcal{A}_S}{4G} + \frac{\mathcal{A}_\infty}{4G} = 4\pi G\left[M_\infty^2 - M_{\rm BH}^2\right]. \tag{22}$$

A nice and useful interpretation of this formula may be obtained by expanding at the leading order $M_\infty = M_{\rm BH} + \delta M_\infty$ (we have checked this approximation to be valid in the BH mass range we are concerned with). One obtains

$$B \simeq 8\pi G M_{\rm BH}\delta M_\infty \simeq \frac{E_{\rm sph}(M_{\rm BH})}{T_{\rm H}}, \tag{23}$$

where in the last passage we have recognised the BH temperature $T_{\rm H} = 1/8\pi G M_{\rm BH}$. The expression (23) tells us that the exponential factor $\exp(-B)$ for the sphaleron transition may be thought of as the standard Boltzmann suppression factor in a thermal environment where the temperature of the system is indeed the Hawking temperature. This interpretation allows as well to smoothly interpolate between the zero and the finite temperature limits. In the case in which the BH is immersed into a plasma at finite temperature T, as in the case of the Primordial BHs (PBHs), the sphaleron baryon number violating rate is expected to go as $\exp(-E_{\rm sph}(M_{\rm BH})/T)$ for $T \gtrsim T_{\rm H}$, fitting the exponential factor (21) at $T \sim T_{\rm H}$. In such limit one must use the VEV of the Higgs field at finite temperature $v(T)$, see Refs. [25,26].

In Fig. 2 we have plotted, on the left panel, the bounce factor B as a function of the BH mass and for two choices of the plasma temperature. One can see that, for enough small BH masses, the bounce B can be much smaller than the vacuum bounce $4\pi/\alpha_W \sim$

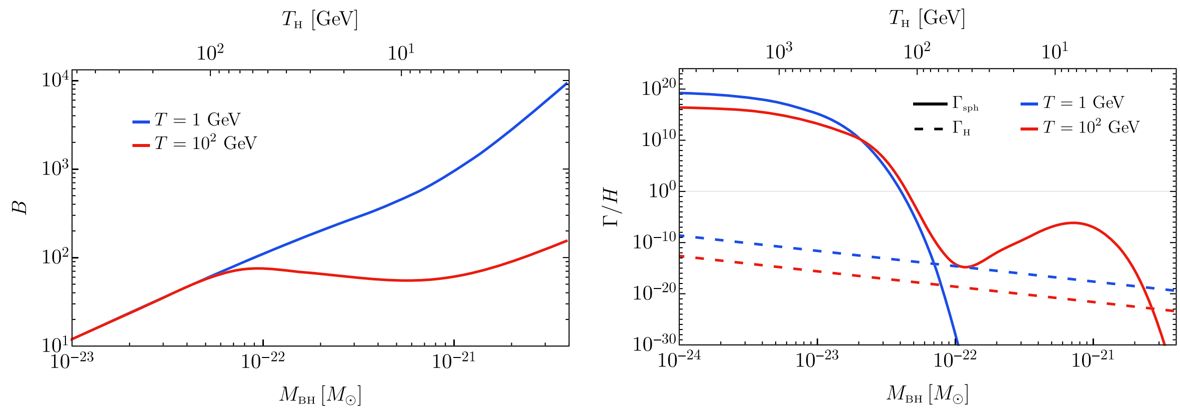

Fig. 2. *Left*: Bounce exponent in terms of the BH mass (and corresponding Hawking temperature) for different temperatures of the plasma. *Right*: Ratio between the sphaleron rate in the presence of a BH of mass $M_{\rm BH}$ and Hubble rate. For comparison, we also show the evaporation rate. Notice that for small masses the sphaleron rate decreases as $\sqrt{B}$, while for larger masses it has a bump due to the plasma temperature effect.

150, reaching values of order unity for masses $M_{\rm BH} \sim 10^{-24} M_\odot$, below which the validity of the computation breaks down. This happens in the region where the expression (23) applies. Moreover, as the BH masses increases, the finite temperature of the thermal bath dominates over the Hawking temperature, leading to a suppression of the bounce exponent.

4. Some further considerations

PBHs with masses of the order of $10^{-22} M_\odot$ may have populated the early Universe, even if with an abundance, normalized to the dark matter one, $f_{\rm PBH} = \Omega_{\rm PBH}/\Omega_{\rm DM} \lesssim 10^{-4}$ to avoid bounds from Big Bang nucleosynthesis [27,28].

If they are formed by the collapse of large overdensities within a horizon volume, their formation temperature is [29]

$$T_{\rm f} \simeq 10^{10} \left(\frac{M_{\rm BH}}{5 \cdot 10^{-22} M_\odot} \right)^{-1/2} \text{GeV}, \tag{24}$$

while their Hawking temperature is in Eq. (13). The light PBHs we are concerned with are always born with a Hawking temperature which is smaller than the plasma temperature. The rate of evaporation for the masses under consideration is quite small and given by [30]

$$\Gamma_{\rm H} \sim 4 \cdot 10^{-33} \left(\frac{10^{-22} M_\odot}{M_{\rm BH}} \right)^{-3} \text{GeV}. \tag{25}$$

In first approximation, one may consider the PBH masses as constant in time for our considerations. A comparison between the baryon number violation rate and the evaporation rate, both in terms of the Hubble rate, can be found in the right panel of Fig. 2. The evaporation rate becomes relevant only for BH masses smaller than $10^{-28} M_\odot$, for which evaporation is effective at temperatures around 100 GeV.

Now, at very high temperatures thermal fluctuations induce unsuppressed baryon number violation through sphaleron transitions till the electroweak phase transition takes place [3]. In the SM this happens at $T_{\rm EW} \simeq 163$ GeV for the current mass of the Higgs. At smaller temperatures and away from the PBHs, the sphalerons are inactive and baryon number violation is suppressed by the exponential $\exp(-E_{\rm sph}(0)/T)$. However, even after the electroweak phase transition, baryon number violation can take place at a rate faster than the rate of the expansion of the universe around the PBHs, see Fig. 2 right panel, where for each BH mass we have taken the maximum between the plasma temperature and the Hawking temperature to evaluate the suppression factor.

Does this represent a threat for the scenarios where the baryon asymmetry of the universe is generated before or at the electroweak phase transition? At the time of formation, the fraction of PBHs per horizon is given by [29]

$$\beta(T_{\rm f}) \simeq 10^{-19} \left(\frac{M_{\rm BH}}{5 \cdot 10^{-22} M_\odot} \right)^{1/2} f_{\rm PBH}. \tag{26}$$

Big Bang nucleosynthesis bounds limit the PBH mass fraction at formation to be $\beta(T_{\rm f}) \lesssim 10^{-23}$ for the range of masses of interest [27]. The number $\mathcal{N}$ of causally independent regions at a time during the radiation-dominated era with temperature T and currently within our horizon is given by $\mathcal{N} \sim 10^{34} (T/\text{GeV})^3$. This means that the number density of PBHs at a given temperature T normalized to the photon number density n_γ is approximately given by

$$\frac{n_{\rm PBH}}{n_\gamma} \frac{1}{\eta} \sim 10^{-34} \left(\frac{\eta}{10^{-9}} \right)^{-1} \left(\frac{M_{\rm BH}}{5 \cdot 10^{-22} M_\odot} \right)^{-1} f_{\rm PBH}, \tag{27}$$

where we have introduced the baryon asymmetry $\eta = n_{\rm b}/n_\gamma$ normalised to the current constrained value [25]. Luckily, the PBH density is too small to have any impact on the pre-existing baryon asymmetry. We believe that such conclusion would be hardly changed envisaging other scenarios of PBH formation in the early universe.

5. Conclusions

We have studied the violation of the baryon number within the SM induced by sphaleron transitions around a BH. Our findings indicate that the bounce for such transitions may be much smaller than the one in the absence of BHs if their Schwarzschild radius is of the order of the electroweak scale. Around PBHs the violation of the baryon number takes place at temperatures below the electroweak phase transition. However our findings indicate that the baryon asymmetry of the universe is unlikely to be wiped out by the presence of PBHs acting as seeds of the sphaleron transitions.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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References

[1] R.F. Klinkhamer, N.S. Manton, Phys. Rev. D 30 (1984) 2212.
[2] G. 't Hooft, Phys. Rev. Lett. 37 (1976) 37.
[3] A. Riotto, M. Trodden, Annu. Rev. Nucl. Part. Sci. 49 (1999) 35–75, arXiv:hep-ph/9901362.
[4] T. Banks, N. Seiberg, Phys. Rev. D 83 (2011) 084019, arXiv:1011.5120 [hep-th].
[5] P. Burda, R. Gregory, I. Moss, Phys. Rev. Lett. 115 (2015) 071303, arXiv:1501.04937 [hep-th].
[6] P. Burda, R. Gregory, I. Moss, J. High Energy Phys. 08 (2015) 114, arXiv:1503.07331 [hep-th].
[7] R. Gregory, K.M. Marshall, F. Michel, I.G. Moss, Phys. Rev. D 98 (8) (2018) 085017, arXiv:1808.02305 [hep-th].
[8] R. Gregory, I.G. Moss, N. Oshita, S. Patrick, J. High Energy Phys. 09 (2020) 135, arXiv:2007.11428 [hep-th].
[9] N. Tetradis, J. Cosmol. Astropart. Phys. 09 (2016) 036, arXiv:1606.04018 [hep-ph].
[10] D. Gorbunov, D. Levkov, A. Panin, J. Cosmol. Astropart. Phys. 10 (2017) 016, arXiv:1704.05399 [astro-ph.CO].
[11] D. Canko, I. Gialamas, G. Jelic-Cizmek, A. Riotto, N. Tetradis, Eur. Phys. J. C 78 (4) (2018) 328, arXiv:1706.01364 [hep-th].
[12] K. Mukaida, M. Yamada, Phys. Rev. D 96 (10) (2017) 103514, arXiv:1706.04523 [hep-th].
[13] K. Kohri, H. Matsui, Phys. Rev. D 98 (12) (2018) 123509, arXiv:1708.02138 [hep-ph].
[14] D.C. Dai, R. Gregory, D. Stojkovic, Phys. Rev. D 101 (12) (2020) 125012, arXiv:1909.00773 [hep-ph].
[15] H. Luckock, I. Moss, Phys. Lett. B 176 (1986) 341–345.
[16] I.G. Moss, N. Shiiki, E. Winstanley, Class. Quantum Gravity 17 (2000) 4161–4174, arXiv:gr-qc/0005007.
[17] T. Hayashi, K. Kamada, N. Oshita, J. Yokoyama, J. High Energy Phys. 08 (2020) 088, arXiv:2005.12808 [hep-th].
[18] P. Candelas, Phys. Rev. D 21 (1980) 2185.
[19] I.G. Moss, Phys. Rev. D 32 (1985) 1333.
[20] S.W. Hawking, Erratum, Commun. Math. Phys. 43 (1975) 199–220, Commun. Math. Phys. 46 (1976) 206.
[21] W.G. Unruh, Phys. Rev. D 14 (1976) 870.
[22] J.B. Hartle, S.W. Hawking, Phys. Rev. D 13 (1976) 2188–2203.
[23] D. Buttazzo, G. Degrassi, P.P. Giardino, G.F. Giudice, F. Sala, A. Salvio, A. Strumia, J. High Energy Phys. 12 (2013) 089, arXiv:1307.3536 [hep-ph].
[24] R. Gregory, I.G. Moss, B. Withers, J. High Energy Phys. 03 (2014) 081, arXiv:1401.0017 [hep-th].
[25] A. Riotto, arXiv:hep-ph/9807454.
[26] M. Quiros, arXiv:hep-ph/9901312.
[27] A.S. Josan, A.M. Green, K.A. Malik, Phys. Rev. D 79 (2009) 103520, arXiv:0903.3184 [astro-ph.CO].
[28] B. Carr, K. Kohri, Y. Sendouda, J. Yokoyama, arXiv:2002.12778 [astro-ph.CO].
[29] For a review, see M. Sasaki, T. Suyama, T. Tanaka, S. Yokoyama, Class. Quantum Gravity 35 (6) (2018) 063001, arXiv:1801.05235 [astro-ph.CO].
[30] D.N. Page, Phys. Rev. D 13 (1976) 198–206.