Matter–antimatter asymmetry induced by a running vacuum coupling

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Received: 7 September 2017 / Accepted: 28 November 2017 / Published online: 11 December 2017

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Abstract We show that a CP-violating interaction induced by a derivative coupling between the running vacuum and a non-conserving baryon current may dynamically break CPT and trigger baryogenesis through an effective chemical potential. By assuming a non-singular class of running vacuum cosmologies which provides a complete cosmic history (from an early inflationary de Sitter stage to the present day quasi-de Sitter acceleration), it is found that an acceptable baryon asymmetry is generated for many different choices of the model parameters. It is interesting that the same ingredient (running vacuum energy density) addresses several open cosmological questions/problems: avoids the initial singularity, provides a smooth exit for primordial inflation, alleviates both the coincidence and the cosmological constant problems, and, finally, is also capable of explaining the generation of matter-antimatter asymmetry in the very early Universe.

1 Introduction

There is a growing body of work on running vacuum cosmologies or an effective dynamical Λ -term [1–6]. These models are motivated by the cosmological constant problem (CCP)—the absence of a satisfactory mechanism whereby the vacuum energy density from all fields can be canceled (or almost canceled) by the fixed cosmological constant. In other words, we do not know how the current, extremely small value of the effective vacuum energy density can be predicted from first principles.

Although successful, the current cosmic concordance model (Λ_0 CDM + Inflation) can also be seen as a collection of many different ingredients that are brought together to explain the complete cosmic evolution over all time

and length scales. Such disparate ingredients are needed in order to smoothly connect the early and late time accelerating regimes. Beyond the enigmatic CCP [7,8], there are also specific challenges for the cosmic concordance model: the initial singularity, the ‘graceful’ exit problem for some popular models of inflation, cosmic coincidence [9] (see also [10–14] for these and other potential problems).

The interest in running vacuum cosmologies can be justified both phenomenologically and from first principles based on quantum field theory in curved space-times. The geometric λ -term (proportional to the metric and usually present on the *l.h.s.* of Einstein’s equations) can be incorporated on the *r.h.s.* of Einstein’s equations as part of the effective energy-momentum tensor (EMT) involving also additional vacuum contributions of the matter sector, $T_{\mu\nu} = \Lambda_E M_{\text{Pl}}^2 g_{\mu\nu}$, where $M_{\text{Pl}} = (8\pi G)^{-1/2} \sim 2.4 \times 10^{18}$ GeV is the reduced Planck mass and Λ_E is the effective Λ -term. The question arises as to what is the total expression of Λ_E ? The vacuum state of all existing fields can be represented by an EMT which reflects the Lorentz invariance of its energy density and pressure. By averaging over all fields the minimally coupled EMT in general relativity reads $\langle T_{\mu\nu} \rangle = \langle \rho \rangle g_{\mu\nu}$. When the geometric λ -term is transferred to the *r.h.s.* of Einstein’s equations one may see that both contributions add, yielding [7,8]

$$\Lambda_E = \lambda + \frac{\langle \rho \rangle}{M_{\text{Pl}}^2}, \quad (1)$$

with the field equations becoming $G^{\mu\nu} = M_{\text{Pl}}^{-2} [T_{(\text{f})}^{\mu\nu} + \Lambda_E M_{\text{Pl}}^2 g^{\mu\nu}]$. Since the Einstein tensor is divergenceless ($\nabla_\mu G^{\mu\nu} = 0$), it is usually argued (even in textbooks) that the vacuum energy density remains constant (*i.e.* $\nabla_\mu \Lambda_E \equiv \partial_\mu \Lambda_E = 0$). However, this result holds only if the cosmic fluid EMT is always separately conserved ($\nabla_\mu T_{(\text{f})}^{\mu\nu} = 0$). In the general case one finds

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$$\nabla_\mu G^{\mu\nu} = 0 \Rightarrow \nabla_\mu T_{(f)}^{\mu\nu} = -g^{\mu\nu} \nabla_\mu \Lambda_E, \quad (2)$$

which implies that a running Λ_E can transfer energy to the cosmic fluid ($T_{(f)}^{\mu\nu}$). In particular, for spatially inhomogeneous and time-dependent space-times even a local dependence of the effective cosmological term, $\Lambda_E \equiv \Lambda_E(x^\mu)$, is possible because energy transfer at different points of the space-time may occur at different rates.

In cosmology, the assumption of a running vacuum energy density is also physically more appealing than the standard view of a fixed Λ_E -term. Since the physical quantities at the background level depend only on time, the condition $\partial_0 \Lambda_E \neq 0$ means that the extremely low value currently observed ($\Lambda_E(t_0) \equiv \Lambda_0$) can be theoretically accommodated. For an aged Universe, the present smallness of $\Lambda_E(t)$ can be seen as a natural consequence of its oldness.

This contemporary view has observational support. Recent work on a class of running vacuum models has shown that a battery of tests including SNIa, BAO, $H(z)$, LSS, BBN, CMB provide a quality fit that is significantly better than the Λ_0 CDM [2–4]. Several authors, following different lines of inquiry, have investigated how the running vacuum transfers energy to the material components. Probably, the most successful path comes from QFT techniques based on the renormalization group (RG) approach in curved space-times. Broadly speaking, one may think that the vacuum energy density ‘runs’ because the effective action inherits quantum effects from the matter sector. Generically, the RG techniques in curved space-times leads to a dependence $\rho_v(H, \dot{H}) = M_{\text{pl}}^2 \Lambda_E(H, \dot{H})$ where H is the Hubble parameter (see [4, 15, 16] and the references therein).

Given the possibility of a running vacuum cosmology, it is also essential to investigate how the *baryogenesis problem*, i.e. the observed matter-(anti)matter asymmetry, can be handled in this context. This old problem becomes more intriguing in the light of the novel running vacuum scenario where the space-time emerges from a pure non-singular de Sitter stage smoothly evolving into radiation, matter and Λ_0 CDM, thereby avoiding the initial singularity and providing a complete description of the expansion history [17, 18]. One may ask whether the baryogenesis mystery can also be explained by the same running vacuum driving the cosmic dynamics.

In this work we address this challenging problem by proposing a new mechanism based on a derivative CP-violating coupling between the running vacuum and the baryon number current. As we shall see, the predicted matter-(anti)matter asymmetry is in good agreement with the observations for many different choices of the relevant cosmic energy scales.

2 Running vacuum baryogenesis

Currently, the baryon asymmetry (B-asymmetry for short) is numerically characterized by the dimensionless η -parameter:

$$5.7 \times 10^{-10} < \eta \equiv \frac{n_B - n_{\bar{B}}}{s} < 6.7 \times 10^{-10}, \quad (3)$$

where n_B , $n_{\bar{B}}$ are the number densities of baryons (anti-baryons, respectively), and s is the present radiation entropy density. The lower and upper limits on the η -parameter come from precision measurements of the primordial deuterium abundance and cosmic background radiation (CMB) acoustic peaks [19–21]. The exact nature of the baryogenesis mechanism is not known yet. Some models have been proposed to explain the B-asymmetry, but there is no consensus in the literature concerning the correct approach [22–26] (see also [27–29] for reviews).

We now discuss a new mechanism appropriate for running $\Lambda_E(t)$ -models (to simplify notation, $\Lambda_E(t) \equiv \Lambda(t)$). The basic ingredient is a derivative coupling between the running vacuum and the baryon current J_B^μ :

$$S = \frac{1}{M_\Lambda^2} \int d^4x \sqrt{-g} (\partial_\mu \Lambda) J_B^\mu, \quad (4)$$

where M_Λ is an unknown cut-off mass scale of the effective decaying vacuum theory whose natural upper bound is the reduced Planck mass M_{pl} . Like the effective Λ -term from (1), the coupling (4) can be viewed either as a particle physics motivated mechanism as in the spontaneous baryogenesis model of [22] which has a coupling of the form

$$S = \frac{1}{M_{\text{SB}}^2} \int d^4x \sqrt{-g} (\partial_\mu \phi) J_B^\mu \quad (5)$$

(where M_{SB} is the cut-off mass scale and ϕ is a scalar field), or a space-time inspired mechanism, as in the gravitational baryogenesis model of [26], which has a coupling of the form

$$S_{\text{GB}} = \frac{1}{M_*^2} \int d^4x \sqrt{-g} (\partial_\mu R) J_B^\mu \quad (6)$$

(where M_* is the cut-off mass scale and R is the Ricci scalar). The proposed coupling in (4) is a mixture of spontaneous baryogenesis and gravitational baryogenesis, keeping the good features of each without some of the drawbacks that both the spontaneous baryogenesis and the gravitational baryogenesis mechanisms have.

As in the gravitational baryogenesis scenario, the multiplicative cut-off mass scale, M_Λ , in (4) has a mass dimension of $[\text{mass}]^{-2}$ instead of $[\text{mass}]^{-1}$ as in the Cohen–Kaplan spontaneous baryogenesis model [22] or similar scenarios [23–25]. The coupling given in (4) shares with the gravitational baryogenesis model several advantages with respect to

the spontaneous baryogenesis model: (1) the scalar field of spontaneous baryogenesis must be assumed to evolve homogeneously in the spatial directions in order to produce the baryon asymmetry while in the present model and that of [26] the homogeneity is naturally built into the model; (2) in the spontaneous baryogenesis model the oscillations of the scalar field tend to time average to zero, which washes out the baryon asymmetry.

The coupling in (4), though different, has advantages to the gravitational baryogenesis model as given by the coupling in (6). Taking into account the spatial homogeneity of the FRW metric one finds that in (6) $\partial_\mu R \rightarrow \dot{R} = 0$ at tree-level. Thus at tree-level the gravitational baryogenesis mechanism does not produce a baryon asymmetry. This is handled in [26] by arguing that to the one-loop level in an $SU(N)$ gauge theory which some fixed number of flavors of quarks one finds that $\dot{R} \neq 0$. However, this assumes that nothing else comes into these one-loop corrections as one goes to higher energy scales. In the model given by the coupling in (4) one automatically gets baryogenesis as soon as one considers a non-constant Λ , without the need to consider higher order loop corrections as in [26].

By requiring a B-violation process in such thermal quasi-equilibrium state it is easy to show that an acceptable B-asymmetry can easily be generated. This happens because the above interaction also implies that CPT is dynamically broken thereby driving the energetics to cause the B-asymmetry. The process is actually very similar to gravitational baryogenesis [26] which relies on the derivative coupling between the Ricci curvature scalar R and the B-current. However, the running vacuum process proposed here has two new interesting features: (1) the same ingredient driving the early accelerating phase of the Universe (i.e. the running vacuum) may also control the baryogenesis process; (2) the running vacuum is always accompanied by particle production and entropy generation [6, 31, 32]. This nonisentropic process is an extra source of T-violation (beyond the freeze out of the B-operator) which as first emphasized by Sakharov [30] is a basic ingredient of baryogenesis.

As remarked before, all quantities, including the running Λ , are only time-dependent (or more accurately depend on the standard FRW time parameter and time foliation). Hence, in the effective Lagrangian (4), one may replace $\partial_\mu \Lambda \rightarrow \partial_0 \Lambda \equiv \dot{\Lambda} < 0$, thereby giving rise to an effective chemical potential. This inequality sign is required by the second law of thermodynamics when the vacuum decays into massless particles in the very early Universe [31, 32]. For a species of particle i carrying a baryon charge of q_i we can define an effective chemical potential as

$$\mu_{Bi} = -q_i \frac{\dot{\Lambda}}{M_\Lambda^2} = \pm \frac{\dot{\Lambda}}{M_\Lambda^2}, \quad (7)$$

where in the last step was assumed that all baryons-(anti)baryons are characterized by baryon numbers ± 1 , respectively. The B-asymmetry produced by the above chemical potential can be written as:

$$\eta = \frac{n_B}{s} \approx -\frac{\dot{\Lambda}}{M_\Lambda^2 T} \Big|_{T=T_D}, \quad (8)$$

with the temperature evaluated at $T = T_D$ when the B-violation mechanism decouples. To obtain the above expression one needs the quasi-equilibrium relationship for the entropy per unit volume when the vacuum decays “adiabatically” [32], $s = \frac{\rho+p}{T} \propto T^3$. Replacing $\dot{\Lambda}$ by $\dot{R}$, where R is the Ricci scalar, (8) becomes exactly the prediction of the gravitational baryogenesis process [26, 33]. In (8) we have approximated the multiplicative factor, $15g_b/4\pi^2 g_* \sim \mathcal{O}(1)$, where $g_b(g_*)$ counts the baryonic (all) degrees of freedom.

The dependence of the η -parameter on the rate of the running Λ , as given by (8), is a key result of this work. In order to illustrate its consequences, we will estimate the generation of the B-asymmetry for a class of non-singular running vacuum cosmology recently proposed [2, 3, 17, 18].

3 Running vacuum cosmology

Let us now consider the dynamic description of decaying $\Lambda(H)$ -models in the context of the flat Friedmann–Robertson–Walker (FRW) type geometry ($c = 1$):

$$ds^2 = dt^2 - a^2(t)(dx^2 + dy^2 + dz^2). \quad (9)$$

In such a background the Einstein field equations for the nonvacuum component plus a cosmological assumes the following form:

$$\rho + M_{Pl}^2 \Lambda(H) = 3M_{Pl}^2 H^2, \quad (10)$$

$$p - M_{Pl}^2 \Lambda(H) = -M_{Pl}^2 (2\dot{H} - 3H^2), \quad (11)$$

where ρ and p are the energy density and pressure of the medium, respectively.

As we shall see, the running vacuum cosmologies adopted here describe a complete cosmological history evolving from an early to a late-time de Sitter state. The vacuum energy density is characterized by a truncated power-series in the Hubble parameter, $H \equiv \dot{a}/a$, in which its dominant term behaves like $\rho_\Lambda(H) \equiv M_{Pl}^2 \Lambda(H) \propto H^{n+2}$, $n > 0$. Such models have some interesting features, among them: (1) regardless the value of n , the Universe starts from a non-singular de Sitter phase ($H = H_I$) and evolves smoothly to the radiation phase, giving a successful model for a graceful exit, (2) the temperature and entropy of radiation is initially zero

(de Sitter vacuum), but the entropy growth due to running $\Lambda(H)$ generates all the present day observed entropy, and the radiation temperature evolution law has also been analytically determined [6, 34], (3) the late-time expansion history is a realistic competitor with the rigid Λ_0 CDM model [2, 3, 6].

The total energy conservation law in running vacuum cosmologies, $u_\alpha T_{(\text{f})}^{\alpha\beta}{}_{;\beta} = -u_\alpha T_{(\Lambda)}^{\alpha\beta}{}_{;\beta}$, reads

$$\dot{\rho} + 3H(\rho + p) = -M_{\text{Pl}}^2 \dot{\Lambda}, \quad (12)$$

where ρ and p are the energy density and pressure of the fluid component. One can see immediately that the above equation can also be directly obtained by combining the cosmological equations (10) and (11) with a variable Λ -term [35–37]. It shows that the variation of Λ is only possible with particle and energy transfer from the running vacuum to the fluid component (see [32] for a general treatment). For simplicity, we consider here a large class of models in which the late-time observed flat Universe is exactly the concordance Λ_0 CDM model while at early times one can have a varying cosmological constant. This kind of model is captured by the phenomenological expression

$$\Lambda(H) = \Lambda_0 + 3\nu H^2 + 3H^2 \left(\frac{H}{H_I} \right)^n, \quad (13)$$

where the power index $n > 0$, ν is some parameter, H_I is an arbitrary inflationary scale and the factors of 3 are introduced for mathematical convenience. Such models were proposed in [6, 17, 18, 38] and were initially purely phenomenologically motivated to give a early inflationary phase (i.e. the last term $\propto H^2 \left(\frac{H}{H_I} \right)^n$), a late-time constant Λ_0 term, and an intermediate term $\propto \nu H^2$. Later models having a varying Λ of the form given in (13) were more rigorously motivated by applying the renormalization group approach from quantum field theory in curved space-times [15, 16]. Note also that for $H \ll H_I$ the model behaves like the cosmic concordance cosmology. From observations over the redshift interval $0 \leq z \leq 1100$ [2–4] one finds that the parameter from the quadratic term, $3\nu H^2$, is $\nu \sim \mathcal{O}(10^{-3})$. Thus at the early times that we consider here and where baryogenesis occurs this term is negligible and we will set $\nu = 0$ in the rest of the paper¹.

Now, by combining the two Friedmann equations above with $\Lambda(H)$ from (13) and taking the limit at early times so

that Λ_0 can be neglected and the model is dominated by relativistic particles, $p_r = \frac{1}{3}\rho_r$, one obtains the evolution equation for the Hubble parameter of the form [38]:

$$\dot{H} + 2H^2 \left(1 - \frac{H^n}{H_I^n} \right) = 0, \quad (14)$$

showing that for $\dot{H} = 0$ one has $H = H_I$, which represents the primordial de Sitter stage. This phase is unstable, and the general analytical solution reads

$$H(a) = \frac{H_I}{[1 + (a/a_{\text{eq}})^{2n}]^{1/n}}, \quad (15)$$

where a_{eq} is the vacuum–radiation equilibrium value of the scale factor, that is, when $\rho_\Lambda = \rho_{\text{rad}}$, a moment coincident with the end of inflation ($\ddot{a} = 0$, $H(a_{\text{eq}}) \equiv H_{\text{end}} = H_I/2^{1/n}$). For $a \ll a_{\text{eq}}$ we find $H = H_I$ (de Sitter) while for $a \gg a_{\text{eq}}$ the solution reduces to the standard radiation phase, $a(t) \propto t^{1/2}$. Therefore, the solution (15) describes a smooth transition from a primeval non-singular Sitter stage to the standard FRW phase regardless of the power index n . This result points to some universality of the process and also suggests a natural “graceful” exit from the early unstable de Sitter stage to the standard radiation epoch when the particle production ends [6, 37, 38].

Under certain conditions, it is well known that the running vacuum may decay into massless particles preserving some equilibrium relations [32]. This happens when the vacuum decays adiabatically. Physically, this means that radiation entropy is generated, effectively massless particles–(anti)particles are created, but the specific entropy per particle remains constant. In this case, it has also been shown that $\rho_{\text{rad}} \propto T_{\text{rad}}^4$ and $n_{\text{rad}} \propto T^3$ [6, 32, 34]. Further, by combining (13) with the first Friedmann equation (10), and taking $(\frac{90}{\pi^2 g_*})^{1/4} \sim 10^{-1/4}$, we obtain the radiation temperature law as a function of H :

$$T_{\text{rad}}(H) = 2^{\frac{n+2}{4n}} T_{\text{end}} \left(\frac{H}{H_I} \right)^{1/2} \left[1 - \left(\frac{H}{H_I} \right)^n \right]^{\frac{1}{4}}, \quad (16)$$

where $T_{\text{end}} = T(H_{\text{end}}) = 10^{-\frac{1}{4}} \sqrt{M_{\text{Pl}} H_I}$ is the temperature at the end of inflation. Note also that $T_{\text{rad}} = 0$ for $H = H_I$. Hence, there is no thermal bath in the begin of the cosmological evolution.

4 Baryogenesis constraints

In order to obtain the strength of the B-asymmetry we need to calculate $\dot{\Lambda}$ and the temperature $T = T_D$ when the B-violation mechanism decouples (see Eq. (8)). Using (13) and

¹ The main effect of ν in the solution (15) for $H(a)$ is to change the energy scale H_I by an effective scale, $\tilde{H}_I = (1 - \nu)^{1/n} H_I$, and the power n is replaced by $\tilde{n} = n(1 - \nu)$ (see Eq. (3.6) in Ref. [5] and subsequent discussion). For the current accepted ν values such corrections are small at the baryogenesis epoch and have been neglected in the present analysis. Naturally, this does not mean that such a quadratic ν -term is negligible at late times (radiation and vacuum-matter stages)

(14) we obtain $\dot{\Lambda}$:

$$\dot{\Lambda} = -6(n+2)H^3 \left(\frac{H}{H_I} \right)^n \left[1 - \frac{H^n}{H_I^n} \right]. \quad (17)$$

To calculate the decoupling temperature T_D we need to specify the B-violating operator to determine when the B-violating process decouples. To simplify the discussion we take a generic GUT, B-violating operator having mass dimension 6 e.g. $\mathcal{L}_{\Delta B} = \frac{1}{M_X^2} \epsilon^{\alpha\beta\gamma} \epsilon^{ab} (\bar{u}_\gamma^c \gamma^\mu q_{\beta a}) (\bar{d}_\alpha^c \gamma_\mu l_b)$ where u, d, q, l stand for up-quark, down-quark, quark and lepton, respectively. The rate of this generic dimension 6 interaction is $\Gamma_{\Delta B} \sim \frac{T_D^5}{M_X^4}$ with M_X being the mass scale connected with $\mathcal{L}_{\Delta B}$. The decoupling temperature is fixed by requiring $\frac{T_D^5}{M_X^4} \sim \Gamma_{\Delta B} \leq H \approx \frac{T_D^2}{M_{\text{Pl}}}$ (the expression for H comes from (16) assuming $H < H_I$ so that the last term is of order 1; also we have taken $2^{-(n+2)/2n} \sqrt{10}$ of order 1). Using this we find the decoupling temperature and decoupling H as

$$T_D \sim \frac{M_X^{4/3}}{M_{\text{Pl}}^{1/3}} \quad \text{and} \quad H_D \sim \frac{M_X^{8/3}}{M_{\text{Pl}}^{5/3}}. \quad (18)$$

Using H_D and T_D and Eqs. (17) and (8) we find

$$\eta \approx 6(n+2) \left(\frac{M_{\text{Pl}}^2}{M_X^2} \right) \left(\frac{H_I}{M_{\text{Pl}}} \right)^{(5n+20)/3} \left(1 - \frac{H_I^{5n/3}}{M_{\text{Pl}}^{5n/3}} \right). \quad (19)$$

For simplicity we have assumed that the inflation scale and B-violating operator mass scale are of the same order – $H_I \sim M_X$. This is in accord with some simple models of inflation where the inflation scale and GUT scale coincide.

The general predictions for η depend on the values assigned to n , M_{Pl}/M_Λ , and H_I/M_{Pl} . These parameters should be chosen to give the observed η and be in agreement with the inflationary scale. In Table 1, we display the running vacuum predictions of the B-asymmetry for a broad set of selected values of the three free parameters. Note as the index n increase M_{Pl}/M_Λ increases if we keep H_I/M_{Pl} in a range that gives $H_I \sim 10^{15} - 10^{16}$ GeV i.e. the GUT scale which is then consistent with our simplifying assumption of $H_I \sim M_X$. The main conclusion from Table 1 is that one can obtain values of η in running vacuum models of the kind discussed here for a range of the relevant sub-Planckian parameters. Thus, this framework provides a successful baryogenesis mechanism driven by a running vacuum model as described here. Finally we note that in all cases $H_I < M_{\text{Pl}}$ thereby showing that trans-Planckian problems are absent.

Table 1 Baryogenesis prediction in a class of running vacuum cosmologies. Note that the above values, which are in good agreement with the current constraints on η – see Eq. (3) – and the inflationary scale [39], are obtained for a range of the model parameters

$n > 0$	$M_{\text{Pl}}/M_\Lambda \geq 1$	$H_I/M_{\text{Pl}} \leq 1$	η
0.1	3.0×10^5	8.2×10^{-4}	6.4×10^{-10}
1.0	3.4×10^5	2.6×10^{-3}	6.0×10^{-10}
2.0	8.2×10^4	9.0×10^{-3}	5.6×10^{-10}
3.0	2.6×10^6	9.6×10^{-3}	5.8×10^{-10}
4.0	5.4×10^8	7.6×10^{-3}	5.8×10^{-10}
5.0	7.1×10^9	9.1×10^{-3}	6.1×10^{-10}
6.0	8.9×10^{11}	8.2×10^{-3}	6.5×10^{-10}

5 Final remarks

We have investigated the early generation of B-asymmetry driven by a non-singular running vacuum cosmology. It should be stressed that the B-violation as discussed here does not require new ingredients. This happens because the running vacuum is the same quantity driving the evolution of the space-time (see Eqs. (4), (8), (13) and (14)). Interestingly, the adopted running vacuum model is also endowed with some remarkable properties like the avoidance of the initial singularity and a smooth exit from inflation to the standard radiation phase thereby producing all the observed entropy. It also alleviates the cosmic coincidence and Λ problems, and, last but not least, we have found the its decay products, through the derivative coupling with the $\Lambda(t)$ term given in (4), is also capable of generating the observed baryon asymmetry.

It is worth noticing that the new mechanism proposed here is different from spontaneous baryogenesis as discussed by several authors [22–25], as well as from gravitationally induced B-asymmetry powered by the Ricci scalar [26, 33, 40], a Gauss–Bonnet term [41] or other higher order curvature invariants. Probably, the closest approach is the gravitational baryogenesis driven by the Ricci scalar in the presence of a running vacuum [40]. However, even in this case there are significant differences. For example, the Ricci scalar is $R = -M_{\text{Pl}}^{-2} T - 4\Lambda$, where $T \equiv \rho - 3p$ is the trace of the matter energy-momentum tensor. For relativistic massless particles (radiation phase, $T \equiv 0$), one finds $\dot{R} = -4\dot{\Lambda}$. Accordingly, gravitational baryogenesis will predict a value of η four times larger than in the present approach. This occurs because the effective chemical potential (and η) in gravitational baryogenesis are proportional to $\dot{R}$ [26] while for the direct running vacuum coupling these quantities are proportional to $\dot{\Lambda}$ (see Eqs. (7), (8)). In addition, for $T \neq 0$, or equivalently, if the EoS parameter is slightly different from $\omega = 1/3$, as discussed in Ref. [26], then gravitational baryogenesis will have two distinct positive contributions associ-

ated with the pair of physical quantities $(\dot{T}, \dot{\Lambda})$. This means that the result $\eta_{\text{GB}} = 4\eta_{\Lambda}$ is a lower bound.

As can be seen from Table 1, the cut-off mass scale of the effective theory (M_{Λ}) does not need to be tuned to Planck mass in order to generate the observed value $\frac{m_{\text{B}}}{s} \sim 10^{-10}$.

Finally, we stress that running vacuum models are able to address a broader class of cosmological observations (including baryogenesis) within a single mechanism [6, 38]. In the case of the cosmological constant problem, for instance, one may compare the vacuum energy densities in the early de Sitter regime and the present day value. By taking the constraints of H_I from Table 1, we see that such a ratio satisfies the inequality $\rho_{\text{vI}}/\rho_{\text{v0}} \simeq \Lambda_I/\Lambda_0 \geq 10^{108}$.

Acknowledgements JASL was partially supported by CNPq, CAPES (PROCAD 2013) and FAPESP (Brazilian Research Agencies). DS is supported by Grant $\Phi.0755$ in fundamental research in natural sciences by the Ministry of Education and Science of Kazakhstan.

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