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Scale-invariance as the origin of dark radiation?

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ABSTRACT

Recent cosmological data favor R^2-inflation and some amount of non-standard dark radiation in the Universe. We show that a framework of high energy scale invariance can explain these data. The spontaneous breaking of this symmetry provides gravity with the Planck mass and particle physics with the electroweak scale. We found that the corresponding massless Nambu–Goldstone bosons – dilatons – are produced at reheating by the inflaton decay right at the amount needed to explain primordial abundances of light chemical elements and anisotropy of the cosmic microwave background. Then we extended the discussion on the interplay with Higgs-inflation and on general class of inflationary models where dilatons are allowed and may form the dark radiation. As a result we put a lower limit on the reheating temperature in a general scale invariant model of inflation.

1. Introduction

Particle physics teaches us that in a renormalizable theory at high energy only dimensionless couplings are relevant. Thus, the Standard Model (SM) becomes scale-invariant at classical level in this limit. Though quantum corrections generally violate scale invariance, one can speculate that at high energy the model is indeed modified to be scale-invariant, which provided the argument by Bardeen [1] can solve the naturalness problem in the SM Higgs sector (suffered from the quadratically divergent quantum corrections to the Higgs boson mass squared, see e.g. [2]). Then spontaneous breaking of the scale invariance provides low energy particle physics with the only (at the tree level) dimensionful parameter of the SM, that is the value of the electroweak scale $v = 246$ GeV.[1]

The same logic may be applied to gravity. Then at high energy the classically scale invariant gravity action[2] contains both scalar curvature R and dilaton X,

$$S_0 = \int d^4x \sqrt{-g}\frac{1}{2}\left[\beta R^2 + (\partial_\mu X)^2 - \xi X^2 R\right], \tag{1}$$

with dimensionless real parameters $\beta, \xi > 0$. Once the scale invariance breaks, dilaton X gains non-zero vacuum expectation value and the last term in (1) yields the Einstein–Hilbert low-energy action. Dilaton remains massless in perturbation theory, so the scale invariance may be maintained at the quantum level, see e.g. [4]. As the Nambu–Goldstone boson, dilaton couples to other fields via derivative thus avoiding bounds on a fifth force.

Remarkably, with R^2-term in gravity action (1), the early Universe exhibits inflationary stage of expansion suggested by Starobinsky [6]. At this stage the Universe becomes flat, homogeneous and isotropic as we know it today. Also, quantum fluctuations of the responsible for inflation scalar degree of freedom in (1) (inflaton, which is also called *scalaron* in this particular model) transform to the adiabatic perturbations of matter with almost scale-invariant power spectrum. These perturbations are believed to be seeds of large scale structures in the present Universe and they are responsible for the anisotropy of the cosmic microwave background (CMB). With β normalized to the amplitude of CMB anisotropy $\delta T/T \sim 10^{-4}$ and ξX^2 fixed by the Planck mass value in order to produce the usual gravity, the action (1) has no free parameters. Therefore, the inflationary dynamics is completely determined. Interestingly, recent analyses of cosmological data [7,8]

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[1] The dark energy may be understood as either an effective cosmological constant emerging after spontaneous breaking of scale invariance or a special dynamics of dilaton field, see e.g. [5].

[2] Quadratic terms in the Riemann and Ricci tensors generally give rise to ghost-like and other instabilities and are omitted hereafter. Since the physics responsible for violation of the scale invariance is also beyond the scope of this paper, we omit the dilaton potential in Eq. (1) and disregard its impact on the early time cosmology. Note that the absence (smallness) of a scale-invariant quartic term X^4 in (1) may be related to vanishing (tiny) cosmological constant at later stages of the Universe expansion [3].

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favor this prediction over those of many other models of inflation driven by a single scalar field.

One may treat these results as a hint of scale invariance at high energy. Yet the theory we consider apart from the Starobinsky model contains also massless dilaton coupled to gravity through the last term in (1). In this Letter we show that this term is also responsible for the scalaron decays into dilatons at post-inflationary reheating. In the late Universe the massless dilatons affect the Universe expansion. Surprisingly, the relic amount of produced massless dilatons is precisely what we need to explain the additional (to active neutrinos) dark radiation component[3] suggested by the recent analyses of CMB anisotropy data [10,7,8,11,12], and favored by the observation of primordial abundance of light chemical elements [13]. We consider this finding as possibly one more hint of scale invariance at high energy.

To complete the study we then discuss the SM Higgs boson sector in the model following Refs. [14–16] and outline the regions of the model parameter space where the SM Higgs contributes to the inflationary dynamics. Finally, we investigate the dilaton production in a general scale invariant model of a single field inflation and set a lower limit on the reheating temperature from avoiding the dilaton overproduction.

2. Dilaton–scalaron inflation

We start from considering the scale invariant extension of the Starobinsky model with action (1). Following [17] we introduce new scalar fields Λ and $\mathcal{R}$ and find the equivalent form of action (1):

$$S=\int d^4x\sqrt{-g}\left[\frac{1}{2}\left(\beta\mathcal{R}^2+(\partial_\mu X)^2-\xi X^2\mathcal{R}\right)-\Lambda\mathcal{R}+\Lambda R\right]. \quad (2)$$

Integrating out auxiliary field $\mathcal{R}$ (solving the corresponding equation of motions for $\mathcal{R}$) we obtain

$$S=\int d^4x\sqrt{-g}\left[\Lambda R+\frac{1}{2}(\partial_\mu X)^2-\frac{1}{2\beta}\left(\Lambda+\frac{1}{2}\xi X^2\right)^2\right]. \quad (3)$$

Going to the Einstein frame through the conformal transformation $g_{\mu\nu}\to\tilde{g}_{\mu\nu}=\Omega^2 g_{\mu\nu}$ with $\Omega^2=-2\Lambda/M_P^2$, and omitting tildes thereafter (all quantities below are evaluated with metric $\tilde{g}_{\mu\nu}$) we arrive at

$$S=\int d^4x\sqrt{-g}\left[-\frac{M_P^2}{2}R+\frac{6M_P^2}{2\omega^2}\left[(\partial_\mu\omega)^2+(\partial_\mu X)^2\right]\right.$$
$$\left.-\frac{M_P^4}{8\beta}\left(1-\frac{6\xi X^2}{\omega^2}\right)^2\right], \quad (4)$$

here $\omega=\sqrt{6}M_P\Omega$, and the reduced Planck mass M_P is defined through the Newtonian gravitational constant G_N as $1/M_P^2=8\pi G_N$. After changing the variables $\omega=r\sin\theta$, $X=r\cos\theta$ the kinetic term K and potential term V become

$$K=\frac{6M_P^2}{2\sin^2\theta}\left((\partial_\mu\log r)^2+(\partial_\mu\theta)^2\right),$$
$$V=\frac{M_P^4}{8\beta}\left(1-6\xi\cot^2\theta\right)^2, \quad (5)$$

or, casting them in terms of new variables

$$\rho=\sqrt{6}M_P\log\frac{r}{M_P},\qquad f-f_0=\sqrt{6}M_P\log\tan\frac{\theta}{2}, \quad (6)$$

[3] Particular models with massless (Nambu–Goldstone) bosons were considered in literature to address the dark radiation problem, see e.g. [9].

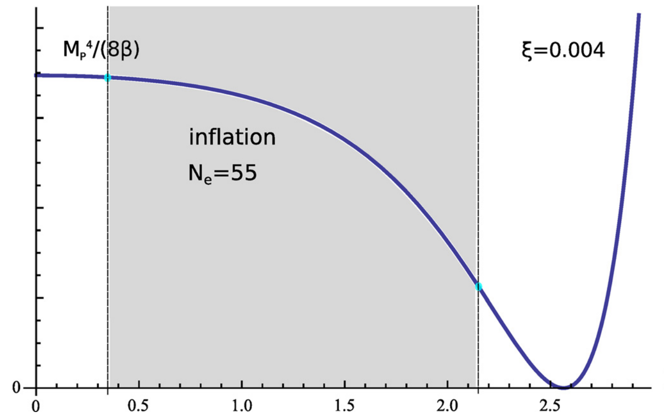

Fig. 1. Inflationary potential: field $F\equiv(f_0-f)/(M_P\sqrt{6})$ slowly moves from close to hilltop $f\simeq f_0$ towards the minimum at $f=0$. The potential is symmetric under reflection $F\to -F$.

we find

$$K=\frac{1}{2}(\partial_\mu\rho)^2\cosh^2\left(\frac{f_0-f}{\sqrt{6}M_P}\right)+\frac{1}{2}(\partial_\mu f)^2,$$
$$V=\frac{M_P^4}{8\beta}\left(1-6\xi\sinh^2\left(\frac{f_0-f}{\sqrt{6}M_P}\right)\right)^2. \quad (7)$$

Both kinetic and potential parts (7) are invariant under reflection $f\to 2f_0-f$. Choosing one of two minima of V to be at $f=0$ implies that integration constant f_0 obeys

$$\sinh^2\left(\frac{f_0}{\sqrt{6}M_P}\right)=\frac{1}{6\xi}. \quad (8)$$

The inflation may occur at values $0<f<f_0$ (or in a mirror interval $f_0<f<2f_0$, that we ignore in what follows), see Fig. 1. The potential is similar to one considered in [14], so for the tilt of scalar perturbations seeded by inflaton fluctuations one has

$$n_s\simeq 1-8\xi\coth(4\xi N_e), \quad (9)$$

where N_e is the number of e-foldings remained till the end of inflation from the moment when perturbations of the CMB-experiments pivot scale $k/a_0=0.002$ Mpc^{-1} exit horizon. To have $N_e\approx 55$ e-folds [18] (since the reheating temperature is about 3.1×10^9 GeV [19] provided scalaron decays to the Higgs bosons) and fit into the favored by cosmological analyses interval $n_s=0.9603\pm 0.0073$ [7], we need

$$\xi<0.004, \quad (10)$$

hence $f_0>6.28\times M_P$. In order to obtain the right value of scalar perturbation amplitude $\Delta\approx 5\times 10^{-5}$ we should choose the parameter β (weakly depending on ξ) to be in the range $(2-0.8)\times 10^9$.

3. Reheating and dilaton production

After inflation the energy is confined in homogeneous oscillating around minimum of the scalaron potential. Scalaron coupling to other fields provides oscillation decay. It reheats the Universe when the Hubble parameter becomes comparable to the inflation decay rate. It is well-known that scalaron couples to any *conformally non-invariant* part of the lagrangian, see discussion in [19, 20]. Within the SM the most relevant is coupling to the Higgs field. Scalaron decay rate to the Higgs bosons is the same as in case of the usual Starobinsky model, and for a more general variant with