Introduction — The idea that Einstein’s gravity may be considered as a large-distance effective theory arising from a spontaneous or dynamical symmetry breakdown in some underlying scale invariant quantum field theory dates back to works of Minkowski [1], Smolin [2], Adler [3], Zee [4], Spokoyny [5], Kleinert and Schmidt [6], and others (see, e.g., Ref. [7] for recent review), even though the incentives can be traced back to 1960’s seminal papers of Zeldovich [8] and Sakharov [9]. Ensuing mechanisms for symmetry breaking are realized typically by spontaneously breaking a scale invariance in appropriate scale-invariant quantum field theory propagating in a curved spacetime [2] or by a Conformal Gravity (CG) which is dynamically broken via additional scalar fields [10, 11].

In particular, the CG has recently attracted renewed attention because local conformal invariance seems to be the key component in number of observationally supported cosmological models. This activity was substantially fueled by Mannheim et al., no-ghosts result [12–15], Smilga’s benign-ghost result [16], and related works on conformal anomaly [17]. The CG has been since revisited from various points of view, e.g., as an alternative to standard Einstein gravity giving a (partial) resolution of a flatness problem [18], or as an explanatory frame for missing matter in galaxies [19] and a possibly vanishing cosmological constant [20]. The CG has also been explored recently in number of theoretical and observational frameworks including conformal supergravity [21], Twistor-String theory [22], asymptotic safety theories [23, 24], black-hole complementarity issue [25], AdS/CFT correspondence [26], and the type Ia supernova (SNIa) and H(z) observational data [27].

Unfortunately, the particle-spectrum of CG does not contain (at least not on-shell) a scalar field. In fact, CG has 6 (on-shell) propagating degrees of freedom; massless spin-2 graviton, massless spin-1 vector boson and massless spin-2 ghost field [21, 28]. Should the Einstein gravity be induced within CG at low energies, the absence of a fundamental scalar poses two imminent problems: a) it is difficult to break a conformal symmetry (either spontaneously or dynamically) without a fundamental spinless boson [21], b) scalar degree of freedom is of a central importance to generate correct primordial density perturbations during inflation [7]. For these reasons an external scalar field is sometimes artificially coupled to CG [10, 11]. In this Letter we wish to point out a subtle fact that a non-dynamical spurion scalar field can be introduced to CG via the Hubbard–Stratonovich transformation without spoiling particle spectrum, (non-perturbative) unitarity, and renormalizability of the CG. The spurion field is actually an imprint of a scalar degree of freedom that would normally be present in the theory if the (local) conformal symmetry would not decouple it.

Quantum Conformal Gravity — CG is a pure metric theory that possesses general coordinate invariance, which augments standard gravity with the additional Weyl symmetry, i.e., invariance under a local rescaling of the metric $g_{\mu\nu}(x) \rightarrow e^{2\alpha(x)}g_{\mu\nu}(x)$, with $\alpha(x)$ being an arbitrary local function. The simplest CG action, i.e., action with both reparametrization and Weyl invariance reads [29, 30]

$$A_{\text{conf}} = -\frac{1}{8\alpha^2} \int d^4x (-g)^{1/2} C_{\lambda\mu\nu \kappa} C^{\lambda\mu\nu \kappa} .$$ (1)
Here $\alpha_c$ is a dimensionless coupling constant and $C_{\lambda\mu\nu\kappa}$ is the Weyl tensor which in 4 space-time dimensions reads

$$C_{\lambda\mu\nu\kappa} = R_{\lambda\mu\nu\kappa} - (g_{\nu\lambda} R_{\mu\kappa} - g_{\kappa\lambda} R_{\mu\nu}) + \frac{1}{3} R g_{\nu\lambda} g_{\mu\kappa},$$

(2)

with $R_{\lambda\mu\nu\kappa}$ being the Riemann curvature tensor, $R_{\lambda\nu} = R_{\lambda\nu}^\mu$ the Ricci tensor, and $R \equiv R^\mu_{\mu}$ the scalar curvature. Throughout we adopt signature $(+,-,-,-)$ and sign conventions of Landau–Lifshitz. With the help of the Gauss–Bonnet theorem one can cast $\mathcal{A}_{\text{conf}}$ into equivalent form (modulo topological term)

$$\mathcal{A}_{\text{conf}} = -\frac{1}{4\alpha_c} \int d^4x (-g)^{1/2} \left[ R_{\mu\nu} R^{\mu\nu} - \frac{1}{3} R^2 \right].$$

(3)

Variation of $\mathcal{A}_{\text{conf}}$ with respect to the metric yields Bach’s field equation [30]

$$2D_\kappa C^{\mu\nu\lambda\kappa} - C^{\mu\lambda\nu\kappa} R_{\lambda\kappa} \equiv B^{\mu\nu} = 0,$$

(4)

where $B^{\mu\nu}$ is the Bach tensor and $D_\alpha$ the Riemannian covariant derivative.

We formally define a quantum field theory of gravity by a functional integral ($\hbar = c = 1$)

$$Z = \int \mathcal{D}g_{\mu\nu} \exp(i\mathcal{A}_{\text{conf}}).$$

(5)

Here $\mathcal{D}g_{\mu\nu}$ denotes the functional-integral measure whose proper treatment involves the Faddeev–Popov gauge fixing of the gauge symmetry $\text{Diff}\times\text{Weyl}(\Sigma_i)$ plus ensuing Fadeev–Popov determinant [31]. Potential local factors $[\det g_{\mu\nu}(x)]^{\omega}$ with Misner’s ($\omega = -5/2$) or De Witt’s ($\omega = (D - 4)/(D + 1)/8$) are omitted in the measure because they do not contribute to the Feynman rules. Their effect is to introduce terms $\omega \delta^4(0) \int \! dx^4 \log(-g)$ into the action, which by Veltman’s rule are set to zero in dimensional regularization. The sum in (5) is a sum over four-topologies, that is, a sum over topologically distinct manifolds $\Sigma_i$ (analogue to the sum over genus in string theory or sum over homotopically inequivalent vacua in the Yang–Mills theory) which can potentially contain topological phase factors, e.g., Euler number of $\Sigma_i$, cf. Refs. [32, 33].

It should be remarked that despite a fourth-order nature of the Bach equation (4) descending from $\mathcal{A}_{\text{conf}}$, it has been recently shown that the would-be ghost states [34] disappear from the energy eigenspectrum and that CG is stable (i.e., unitary) [12, 13]. Also, the conformal instability typical for the Euclidean quantum gravity is not presents in CG.

Most importantly, the quadratic-curvature action [35] is power-counting renormalisable and asymptotically free ($\beta$-function for $\alpha_c$ is negative) [35, 36].

Uncompleting the $R^2$-term — Here we wish to point out that the large number of derivatives in the free graviton propagator implied by [3] makes fluctuations so violent so that the theory might spontaneously create a new mass term. This phenomenon is indeed known to happen in number of higher-derivative systems ranging from biomembranes [36] through string theories with extrinsic curvature [37, 38], to gravity-like theories [39]. For instance, in biomembranes and stiff strings the ensuing mass term can be identified with a tension. We shall now show that an analogous mechanism spontaneously generates the Starobinsky action [40]

$$\mathcal{A}_{\text{St}} = -\frac{1}{2\kappa^2} \int d^4x (-g)^{1/2}(R - \xi^2 R^2).$$

(6)

Here $\kappa^2 = 8\pi G_N$ where $G_N = 1/m_p^2$ is Newton’s (gravitational) constant and $m_p$ is the Planck mass. Starobinsky’s parameter $\xi$ is related to the inflational scale and by the Planck satellite data $\xi/\kappa \sim 10^5$ (cf. Ref. [41]). The minus sign in front of $R^2$-term is a consequence of the Landau–Lifshitz convention [42].

In order to see how the spontaneous generation of $R^2$ comes about we first observe that the $R^2$-part of the action [36] is the global scale-invariant expression. This is because under infinitesimal Weyl transformation $g_{\mu\nu} \rightarrow g_{\mu\nu} + 2\alpha(x) g_{\mu\nu}$ while $R \rightarrow |1 - 2\alpha(x)| R - 6D^2\alpha(x)$ (the covariant derivative $D_\mu$ is with respect to $g_{\mu\nu}$). Since $g \rightarrow (1 + 8\alpha(x)) g$, the $R^2$-term part of the action will be scale invariant provide $D^2\alpha(x) = 0$. The $R^2$-part of the action can be decomposed by using the Hubbard–Stratonovich (HS) transformation [43, 44]

$$\exp(i\mathcal{A}_{\text{ghi}}) \equiv \exp \left( -i \frac{1}{12\alpha_c} \int d^4x (-g)^{1/2} R^2 \right)$$

$$= \int \mathcal{D}\lambda \exp \left[ -i \int d^4x (-g)^{1/2} \left( \frac{3\alpha^2}{4} \lambda^2 + \frac{\lambda R}{2} \right) \right].$$

(7)

Although an auxiliary field $\lambda(x)$ in (7) does not have a bare kinetic term, the conformal symmetry allows to rescale the metric so that a kinetic term can easily be generated. For instance, when $g_{\mu\nu} \rightarrow |\lambda|^{-1} g_{\mu\nu}$ then $\mathcal{A}_{\text{ghi}}$ goes to

$$\int d^4x (-g)^{1/2} \left( \frac{\lambda R}{2|\lambda|} + \frac{3}{4\lambda^2} \partial_{\mu}\lambda \partial^\mu \lambda - \frac{3\alpha^2}{4} \right),$$

(8)

(and other higher-order derivatives of $\lambda$ will come from the remaining $R^{\mu\nu} R_{\mu\nu}$-term). Since the $\lambda$-kinetic term depends on the conformal scaling $\lambda$-kinematics is gauge dependent, implying that $\lambda$ cannot represent a physical field. On the other hand, when the conformal symmetry breaks down then the $\lambda$-field is trapped in a particular (broken) phase with specific kinetic and potential terms. This will be seen shortly.

To proceed, it is helpful to separate the $\lambda$-field into a background field $\lambda$ corresponding to the VEV of $\lambda$
and fluctuations $\delta \lambda$ which have only nonzero momenta. Of course, the fluctuations must be included to make the theory completely equivalent to the original \([5]\). In the following we employ the standard effective-action strategy, i.e., neglect all terms involving $\delta \lambda$, and take the saddle-point approximation to the remaining integral over $\lambda$.

As will be seen shortly, $\lambda$ spontaneously develops a positive VEV, so that the sign of the $R$-term in \([7]\) coincides with the sign of the Einstein term. With the benefit of hindsight we introduce an arbitrary mixing angle $\theta$ and write formally $A_{\text{gsm}} = C^2 A_{\text{gsm}} - S^2 A_{\text{gsm}}$, where $C \equiv \cosh \theta, S \equiv \sin \theta$. Applying the HS-transformation only to the (S$^2$A$^2$gsm)-part we get, after a formal replacement $\alpha^2_c \to -\alpha^2_c / S^2$ in \([7]\)

\[
A_{\text{gsm}} = \frac{C^2}{12 \alpha_c^2} \int \! d^4x (g)^{1/2} R^2 - \frac{1}{2 \kappa^2} \int \! d^4x (g)^{1/2} \lambda R \\
+ \frac{3 \alpha_c^2}{4 S^2 \kappa^4} \int \! d^4x (g)^{-1/2} \lambda^2.
\]  

\[\tag{9}\]

$\mathcal{A}_{\text{conf}} = \frac{1}{16} \int \! d^4x h^{\mu \nu} \Box^2 h_{\mu \nu} - \frac{1}{8} \int \! d^4x \partial_{\lambda} h^{\mu \nu} H_{\mu \nu} \partial_\rho h^{\rho \nu} + \frac{1}{4} \left( \frac{1}{4} - \frac{C^2}{3} \right) \int \! d^4x \bar{h} \Box h \\
+ \frac{1}{2 \kappa^2} \int \! d^4x \left( -\alpha_c \bar{\lambda} \Box h - \frac{\alpha^2_c}{4} \bar{\lambda} \Box h - \frac{\alpha^2_c}{2} \partial_{\lambda} h^{\mu \nu} \bar{\lambda} \Box^{-1} H_{\mu \nu} \partial_\rho h^{\rho \nu} + \frac{\alpha^2_c}{4} h^{\mu \nu} \bar{\lambda} \Box h_{\mu \nu} \right) - \frac{3 \alpha_c^2}{4 S^2 \kappa^4} \int \! d^4x \bar{\lambda}^2 \\
= \int \! d^4x h^{\mu \nu} \mathfrak{A} \Box^2 h_{\mu \nu} + \int \! d^4x \partial_{\lambda} h^{\mu \nu} \mathfrak{B} H_{\mu \nu} \partial_\rho h^{\rho \nu} + \int \! d^4x \bar{h} \Box \Box h + \frac{3 \alpha_c^2}{4 S^2 \kappa^4} \int \! d^4x \bar{\lambda}^2,
\]

\[\mathfrak{A} = \frac{1}{16} + \frac{\alpha^2_c \bar{\lambda} \Box^{-1}}{(8 \kappa^2)}, \quad \mathfrak{B} = -\mathfrak{A}/2, \quad \mathfrak{C} = \frac{1}{4} \left( \frac{1}{4} - \frac{C^2}{3} \right) - \frac{\alpha^2_c \bar{\lambda} \Box^{-1}}{(8 \kappa^2)}, \quad \tag{10}\]

where $H_{\mu \nu} = 1/2 \partial_\rho h^{\rho \nu} - \Box \eta_{\mu \nu}$ and $\bar{h} = h^{\mu \nu} - \Box^{-1} \partial_\rho h^{\rho \nu}$.

A phenomenologically consistent long-range behavior of the gravitational field is ensured if $\bar{\lambda} = 1$. To see that such a solution exists at low enough energies we calculate the one-loop contribution to the Minkowski effective action. This is obtained by functionally integrating out the fields $h_{\mu \nu}$ in the exponential $e^{i \mathcal{A}_{\text{conf}}}$ in which $\lambda$ is approximated by its VEV, i.e., $\lambda$. The result is $e^{-i \Omega_4 V_{\text{eff}}}$, where $\Omega_4$ is the total four-volume of the universe, and $V_{\text{eff}}$ is the effective potential. Form \([10]\) is particularly convenient for the gauge fixings \([10, 21]\): $\chi^\nu = \partial_\mu h^{\mu \nu} = \zeta^\nu (x)$ (coordinate gauge) and $\chi \equiv \bar{h} = \zeta (x)$ (conformal gauge). Here $\zeta^\nu (x)$ and $\zeta (x)$ are arbitrary functions of $x$. Using 't Hooft’s averaging trick \([46]\):

\[
\delta [\chi - \zeta] \to \int \mathcal{D} \zeta e^{i \mathcal{A} \zeta (\det \mathcal{H})^{1/2} \delta [\chi - \zeta]} \\
= e^{i \mathcal{A} \chi \mathcal{H} (\det \mathcal{H})^{1/2}} (\mathcal{H} \text{ is an arbitrary symmetric operator and doing some straightforward computations we obtain the zero-genus (fixed topology) contribution to partition function}
\]

\[
Z_0 = \mathcal{N} (\det \mathcal{M}_{\text{FP}}) (\det H_{\mu \nu} \det (\Box^2) h_{\mu \nu})^{-1/2} [\det (\Box^2) h_{\mu \nu}]^{-1/2} (\det \mathfrak{A})^{1/2} (\det \mathfrak{C})^{-1} - 3 e^{i \Omega_4 3 \alpha^2_c \bar{\lambda}^2 / (4 S^2 \kappa^4)} \\
= \mathcal{N} (\det (-\Box))^{-1/2} (\det \mathfrak{C})^{1/2} (\det \mathfrak{A})^{-1} e^{-i \Omega_4 3 \alpha^2_c \bar{\lambda}^2 / (4 S^2 \kappa^4)}, \quad \tag{12}\]

Here the additional rescaling $\lambda \to \lambda / \kappa^2$ was included because we want our theory to eventually induce Einstein’s action after $\lambda (x)$ acquires a VEV.

Let us now show that the fluctuations of the metric $g_{\mu \nu}$ can achieve this. In particular, we find a set of parameters in the model parameter space for which $\bar{\lambda} = \langle \lambda \rangle = 1$. As a result, the long-range behavior of our theory will coincide with that of Starobinsky’s $f(R)$-model.

\textbf{Emergence of Starobinsky’s model — We proceed by splitting the spacetime metric into the flat Minkowski background plus a fluctuation $h_{\mu \nu}$ defined by $g_{\mu \nu} = \eta_{\mu \nu} + \alpha_c h_{\mu \nu}$ (realizing that $\alpha_c \sim C \kappa / \xi$, and then expanding the Lagrangian in \([5]\) (including the explicit form \([9]\)) to the second order in $\alpha_c$. Omitting total derivatives, using the simple weak-field relations of Appendix A and setting $\lambda = \bar{\lambda}$, we end up with the following outcome (here $\Box \equiv \partial^2$)
\[
(M_{FP})_{\mu\nu} = -\Box h_{\mu\nu} - \partial_{\mu}\partial_{\nu}
\] is the Faddeev–Popov operator for coordinate gauge \([47]\). The factor \([\det(-\Box)]^{1/2}\) correctly indicates that that number of propagating modes in the linearized CG is 6 (cf. Ref. [28]). From (17) the one-loop \(V_{\text{eff}}\) reads
\[
V_{\text{eff}} = \frac{i}{2} \int \frac{d^4k}{(2\pi)^4} \ln \left(k^2 - \frac{6\alpha_2^2\lambda}{\kappa^2(4S^2 + 1)}\right) - \frac{6i}{2} \int \frac{d^4k}{(2\pi)^4} \ln \left(k^2 - \frac{2\alpha_2^2\lambda}{\kappa^2}\right) - \frac{3\alpha^2}{4S^2\kappa^4}\lambda^2.
\]
(13)
The prime indicates a trivial subtraction of the zero-mode. Note that for (assumed) \(\lambda > 0\) the ensuing massive pole is physical only when \(\theta \in (-\arcsin(1/4), \infty)\). The integral over \(k\) can be evaluated, e.g., with the help of dimensional regularization \((D = 4 - 2\epsilon)\) in which case it yields
\[
V_{\text{eff}} = -\frac{9\alpha_2^4\lambda^2}{16\pi^2\kappa^4(4S^2 + 1)^2} \ln \left(\frac{6\alpha_2^2\lambda}{(1 + 4S^2)\kappa^2\mu^2}\right) - \frac{3}{2}
\]
\[+ \frac{3\alpha_2^2\lambda^2}{8\pi^2\kappa^4} \ln \left(\frac{2\alpha_2^2\lambda}{\kappa^2}\right) - \frac{3}{2} - \frac{3\alpha^2}{4S^2\kappa^4}\lambda^2,
\]
(14)
where \(\Lambda = \sqrt{4\pi\mu\epsilon^{1/2}}\), \(\mu\) is an arbitrary renormalization scale and \(\gamma\) is the Euler–Mascheroni constant. To obtain a finite result as \(\epsilon \to 0\) we utilize the \(\overline{MS}\) renormalization scheme. This fixes the counterterm so that
\[
V_{\text{eff}} = -\frac{9\alpha_2^4\lambda^2}{16\pi^2\kappa^4(4S^2 + 1)^2} \ln \left(\frac{6\alpha_2^2\lambda}{(1 + 4S^2)\kappa^2\mu^2}\right) - \frac{3}{2}
\]
\[+ \frac{3\alpha_2^2\lambda^2}{8\pi^2\kappa^4} \ln \left(\frac{2\alpha_2^2\lambda}{\kappa^2}\right) - \frac{3}{2} - \frac{3\alpha^2}{4S^2\kappa^4}\lambda^2,
\]
(15)
with \(\mu^2\) being the subtraction point.
The saddle point in \(\lambda\) corresponding to the VEV is determined by the vanishing of \(V_{\lambda} \equiv \partial V_{\text{eff}}/\partial \lambda\). This yields the minimal \(V_{\text{eff}}\) for
\[
\lambda(S) = \exp \left(\frac{3\alpha_2^2S^2 \ln \left(\frac{3}{4S^2 + 1}\right)}{\alpha^2S^2(32S^4 + 16S^2 - 1)}\right)
\]
\[\times \frac{\kappa^2\mu^2\epsilon}{2\alpha_2^2\epsilon}.
\]
(16)
In this case \(V_{\text{eff}} < 0\) for \(S^2 > (\sqrt{6} - 2)/8 \approx 0.056\), irrespective of actual values of \(\alpha\) and \(\kappa\). A trivial solution of \(V_{\lambda} = 0\), namely \(\lambda(S) = 0\) yields \(V_{\text{eff}} = 0\) and hence it represents a local maximum (i.e., unstable solution) for the above range of \(S^2\).

Although the full theory described by the action \([10]\) is independent of the mixing angle \(\theta\), the truncation of the perturbation series after a finite loop order in the fluctuating \(h_{\mu\nu}\)-field spoils this independence. The optimal result is obtained by utilizing the principle of minimal sensitivity \([38]\) known from the renormalization-group calculus. The principle of minimal sensitivity is at the heart of the \(\delta\)-perturbation expansion \([49]\) and variational perturbation expansion \([51, 52]\). There, if the perturbation theory depends on an unphysical parameter, say \(\theta\), the best result is achieved if each order has the weakest possible dependence on the parameter \(\theta\). Consequently, at the one-loop level the value of \(\theta\) is determined from the vanishing of the derivative of \(V_{\text{eff}}\) with respect to \(S^2\). By setting \(V_{S^2} \equiv \partial V_{\text{eff}}/\partial S^2\), we have
\[
d\frac{dV_{\text{eff}}}{dS^2} = \frac{\partial \lambda(\theta)}{\partial S^2} V_{\lambda} + V_{S^2} = V_{S^2} = 0.
\]
(17)
This is equivalent to the equation
\[
(128S^6 + 96S^4 + 36S^2 - 1) = \frac{12\alpha^2\ln \left(\frac{4S^2 + 1}{3}\right)}{\pi^2(32S^4 + 16S^2 - 1)}(18)
\]
which admits two branches of real solutions; either \(S^2 = 0.0259237 - 0.0000197\alpha^2 + \mathcal{O}(\alpha^4)\) which, however, does not give a stable \(\lambda(S)\) (as \(V_{\text{eff}} > 0\)) or infinite (but \(\lambda(S)\)-stable) which withing the range of validity of our one-loop approximations means that \(S\) is at most \(\sim \xi/\kappa \sim 10^3\). Consequently, only sufficiently large values of \(S\) are optimal. From Eq. (16) we deduce the one-loop VEV (to order \(\mathcal{O}(1/S^4)\))
\[
\lambda = \frac{\kappa^2\mu^2}{2\alpha_2^2\epsilon} e^{1+2\pi^2/\alpha^2S^2} \sim \frac{\kappa^2\mu^2}{2\alpha_2^2\epsilon} e^{1+2\pi^2\kappa^2/\alpha^2\epsilon^2}.
\]
(19)
In particular, for any value of the dimensionless coupling strength \(\alpha_c\), we can choose the renormalization mass scale \(\mu\) in such a way that \(\lambda\) has the value 1, that will guarantee phenomenologically correct gravitational forces at long distances. VEV \(\lambda\) is thus the dimensionally transmuted parameter of the massless CG. Its role here is completely analogous to the role of the dimensionally transmuted coupling constant in the Coleman–Weinberg treatment of the massless scalar electrodynamics \([53]\). Namely, we have traded a dimensionless parameter \(\alpha_c\) for a dimensionfull parameter \(\lambda/\kappa^2\) (which does not exist in the symmetric phase).

By assuming that in the broken phase a cosmologically relevant metric is that of Friedmann–Lamaitre–Robertson–Walker (FRLW), then modulo topological term the additional condition
\[
\int d^4x (-g)^{1/2} 3R_{\mu\nu}R^{\mu\nu} = \int d^4x (-g)^{1/2} R^2,
\]
(20)
holds due to a conformal flatness of the FRLW metric \([54]\). Combining \([10, 109, 110]\), the low-energy limit of \(A_{\text{conf}}\) in the broken phase reads
\[
A_{\text{conf. b.}} = -\frac{1}{2\kappa^2} \int d^4x (-g)^{1/2} (R - \xi^2 R^2 - 2\Lambda),
\]
(21)
with
\[
\xi^2 = \frac{\kappa^2 S^2}{6\alpha_2^2}, \quad \Lambda = \frac{3\alpha_2^2}{4S^2\kappa^4}.
\]
(22)
We should stress that Λ is entirely of a geometric origin (it descends from the CG) and it enters in (21) with the opposite sign in comparison with the matter-sector induced (de Sitter) cosmological constant.

**Gradient term for λ** — The local conformal symmetry dictates that the scalar degree of freedom must decouple from the on-shell spectrum of the CG [21, 28], whereas in theories without conformal invariance (but with the same tensorial content) the scalar field does appear in spectrum [11, 28, 34]. When the conformal symmetry is broken the scalar field reappears through radiatively induced gradient term of the spurion field λ. The explicit form of the kinetic term (namely its overall sign!) can be decided from the momentum-dependent part of the λ self-energy Σλ. This can be streamlined by considering in (10) slowly fluctuating λ instead of fixed ˘λ. Since the lowest-order contribution to Σλ comes from coupling to h, the only relevant substitutions in (10) are: λ□h → λ□h (which stops to be a total-derivative) and ˜h□h → λhναβ□αβ = πµν(λµν)□h (P(0)µν,αβ = πµνπαβ/3 is the spin-0 projection, and πµν = ηµν − δµh−1δν is the transverse vector projection). In the leading αc-order, one can neglect αcδµλ with respect to αcδµλ and complete the square in (10) as follows

\[ -\frac{αcλ}{2κ}□h + ˘h□2h \]

\[ \rightarrow ˘h□2h - \frac{αc2}{16κ2} λ(□−2□c−1)λ \]

\[ \approx ˘h□2h + \frac{1}{2κ2}λ□h. \]  

(23)

The last approximation holds for \( 1 \ll αc2□−1/(Cκ)^2 \sim □−1/κ^2 \sim 10^{28}□−1 \), and thus in the large-scale cosmology where only low-frequency modes of scalar fields (e.g., λ) are observationally relevant. The square completion procedure employed in (23) changes the (conformal) gauge fixing condition, albeit the only effect of this modification is a redefinition of the function ζ.

Because of a minus sign in front of \( A_{\text{conf}} \) in (10), the actual kinetic term is \(-\frac{1}{2κ}λ□h \sim 2κλδµδνλ □h \) which is positive. As a result, λ morphs into a genuine (non-ghost, non-tachyonic) propagating scalar mode.

In passing, we note that since \( V_{\text{eff}} \) in the broken phase is bounded from below and the kinetic energy is positive (i.e., vacuum decay is prevented), the broken one-loop linearized CG does not possess ghost states.

**Cosmological implications** — Recent polarisation data from Planck and WMAP satellites [11] support inflationary models with small tensor-to-scalar ratio: \( r < 0.12 \) at 95% CL. These include, e.g., the Starobinski model [38], the non-minimally coupled model (\( \propto φ^2 R/2 \)) with a \( V(φ) \propto φ^4/4 \)-potential, and an inflation model based on a Higgs field [41]. In the SM the linear Einstein term determines the long-wavelength behavior while the \( R^2 \)-term dominates short distances and drives inflation. In phenomenological cosmology, the SM represents metric gravity with a curvature-driven inflation. In particular, it does not contain any fundamental scalar field that could be an inflaton, even though a scalar field/inflaton formally appears when transforming the SM to the Einstein frame [55].

SM emerges naturally in CG in the weak-field sector of the broken phase where the action \( A_{\text{conf,b,λ}} \) reads

\[ \frac{1}{2κ^2} \int d^4x (-g)^{1/2} \left( λR - ξ^2 R^2 - \frac{(∂µλ)^2}{λ} - 2Λλ^2 \right) \]

\[ \frac{λ→λ}{2κ^2} \int d^4x (-g)^{1/2} (R - ξ^2 R^2 - 2Λ). \]  

(24)

Similarly, as in the usual SM, one can set up for \( A_{\text{conf,b,λ}} \) a dual description in terms of a non-minimally coupled auxiliary scalar field φ with the action [2, 55]

\[ A_{φ,E} = -\frac{1}{κ^2} \int d^4x (-g)^{1/2} \left[ \frac{R}{2} - \frac{3ξ(∂µφ)(∂µφ)}{(λ + 2ξφ)^2} - \frac{3(∂µφ)^2}{2(λ + 2ξφ)^2} - \frac{3(∂µφ)^2}{(λ + 2ξφ)^2} \right. \]

\[ \left. + \frac{φ^2}{2(λ + 2ξφ)^2} - \frac{Λλ^2}{(λ + 2ξφ)^2} \right]. \]  

(26)

The above metric rescalling is valid only for the metric-signature-preserving transformation, i.e., only when \( λ + 2ξφ > 0 \). The action (26) can be brought into a diagonal form if we pass from fields \( \{λ, φ\} \) to \( \{ψ, λ\} \) where the new field ψ is obtained via the redefinition \( φ = \exp(\sqrt{2/3}|ψ|) - λ/(2ξ) \). In terms of ψ the action reads

\[ A_{ψ,E} = -\frac{1}{κ^2} \int d^4x (-g)^{1/2} \left[ \frac{R}{2} - \frac{1}{2}(∂µψ)^2 + U(ψ, λ) \right. \]

\[ -e^{-\sqrt{2/3}|ψ|}(∂µψ)^2 \frac{2λ}{2λ} \left. \right] \],  

(27)

where \( U(ψ, λ) = \frac{1}{8ξ^2} \left( 1 - 2λe^{-\sqrt{2/3}|ψ|} \right) \), with ξ from [22]. The strength of λ-field oscillations is controlled by
the size of a coefficient in front of the \( \lambda \)-gradient term \[50\], i.e., \( e^{-\sqrt{2/3} |\psi|/\kappa^2} \) (more precisely, the local fluctuations square width \( (\lambda(x-\bar{\lambda}))^2 \sim \kappa^2 e^{\sqrt{2/3}|\psi(x)|} \)). At large values of the dimensionless scalar field \( \psi \), i.e., at values of the dimensionful field \( \bar{\psi} = \psi/\kappa \) that are large compared to the Planck scale, the gradient coefficient is very small and \( A \)-field severely fluctuates. Assuming that CG was broken before the onset of inflation, then after a brief period of violent oscillations the \( \lambda \)-fluctuations are strongly damped \[51\] at \( \bar{\psi} \lesssim 10^5 m_p \). From then on, the \( \lambda \)-field settles at its potential minimum at \( \bar{\psi} \equiv \bar{\psi} \) (collective) scalar field \( \psi \) playing the role of inflaton. Using the slow-roll parameters

\[
\epsilon = \frac{1}{2} m_p^2 \left( \frac{\partial \psi U(\psi, \lambda)}{U(\psi, \lambda)} \right)^2, \quad \eta = m_p^2 \frac{\partial^2 U(\psi, \lambda)}{U(\psi, \lambda)}, \quad (28)
\]

(\( \partial_\psi \equiv \partial / \partial \psi \)) one can write down the tensor-to-scalar ratio \( r \) and the spectral index \( n_s \) in the slow-role approximation as \[41\]

\[
r = 16 \epsilon, \quad n_s = 1 - 6 \epsilon + 2 \eta. \quad (29)
\]

In terms of the number \( N \) of \( e \)-folds left to the end of inflation

\[
N = -\kappa^2 \int_{\psi_f}^{\psi_i} d\psi \frac{U(\psi, \lambda)}{\partial_\psi U(\psi, \lambda)} \approx \frac{3}{4 \lambda} e^{\sqrt{2/3} |\psi|}, \quad (30)
\]

(\( \psi_f \) represents the values of the inflaton at the end of inflation, i.e., when \( e^{-\sqrt{2/3} |\psi|} \approx 1 \)) one gets

\[
n_s \approx 1 - \frac{2}{N}, \quad r \approx \frac{12}{N^2}, \quad (31)
\]

which for \( N = 50 - 60 \) (i.e., values relevant for the CMB) is remarkably consistent with Planck data \[41\].

While during the inflation, the \( \lambda \)-field is constant (due to a large coefficient in front of the gradient term) allowing a large-valued inflaton field descend slowly from potential plateau, inflation ends gradually when \( \lambda \) regains its canonical kinetic term, and a small-valued inflaton field picks up kinetic energy. From \[27\] the dominant interaction channels at small \( |\psi| \) is \( (\partial_\psi \lambda)^2 |\psi| \), hence the vacuum energy density stored in the inflaton field is transferred to the \( \lambda \) field via inflaton decay \( \psi \rightarrow \lambda + \lambda \) (reheating), possibly preceded by a non-perturbative stage (preheating).

Note also, that the gravi-cosmological constant \( \Lambda \) that was instrumental in setting the inflaton potential in \[27\] has the opposite sign when compared with ordinary (matter-sector induced) cosmological constant. Since the conformal symmetry prohibits the existence of a (scale-full) cosmological constant, the gravi-cosmological constant must correspond to a scale at which the conformal symmetry breaks, which in turn determines the cut-off scale of the scalaron. The magnitude of \( \xi \) in the SM is closely linked to the scale of inflation \[32\]. Using the values relevant for the CMB with \( 50 - 60 \) \( e \)-foldings, the Planck data \[33\] require \( \xi \sim 10^{-13} \text{GeV} \) or equivalently \( \xi/\kappa \sim 10^5 \). Thus from \[22\] the vacuum energy density is \( \rho_\Lambda \equiv \Lambda/\kappa^2 \sim 10^{-10}(10^{18} \text{GeV})^4 \approx 2 \times 10^{100} \text{erg/cm}^3 \), which corresponds to a zero-point energy density of a scalaron with an ultraviolet cut-off at \( 10^{15} - 10^{16} \text{GeV} \). This is in a range of the GUT inflationary scale. For compatibility with an inflationary-induced large structure formation the conformal symmetry should be broken before (or during) inflation \[34\]. This can be naturally included in a broader theoretical context of “conformal inflation” paradigm, which has been the thrust of much of the recent research \[37\]-\[59\]. Let us also notice that an existence of a single scalar field with cutoff at the GUT scale and coupled to broken CG (e.g., \( \lambda \) or GUT Higgs field) would contribute with a zero-point energy that could substantially reduce or eliminate \( \Lambda \).

**Conclusions** — To conclude, we have shown that a spurion-field mediated spontaneous symmetry breakdown of CG is capable of transforming a purely metric conformal gravity to an effective scalar-tensor gravity. This offers a new paradigm to understand inflationary and large-scale cosmology. In particular, we have shown that the low-energy dynamics in the broken phase is described by a Starobinsky-type \( f(R) \)-model which can be mapped on a two-field hybrid inflationary model. A dimensional transmutation ties up together values of Starobinsky’s inflation parameter \( \xi \) and cosmological constant \( \Lambda \). This fixes the symmetry-breakdown scale for CG to be roughly the GUT inflationary scale. Despite its simplicity, the presented paradigm reproduces not only phenomenologically acceptable picture of the large-scale Universe that is compatible with present Planck and WMAP data but it also provides a viable mechanism for reheating mechanism.

If BICEP2 data that support a large \( r > 0.16 \) (i.e., large-field) inflationary models are confirmed by other subsequent experiments, the conventional Starobinsky-type inflationary potential will be excluded. In turn, this would also invalidate the outlined scenario. Clearly one could present modified scheme in which value of \( r \) would be large, e.g., by considering a tensorial HS transformation for \( R_{\mu\nu}^2 \) or superconformal supergravity. These scenarios seem, however, at present less appealing.

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**Appendix A** — Here we collect some technical points used in the text. The weak-field expansion of \( A_{\text{conf}} \) is
based on the fluctuating field $h_{\mu\nu}$: $g_{\mu\nu} = \eta_{\mu\nu} + \alpha_c h_{\mu\nu}$. This gives $R_{\mu\nu,k} = \frac{\alpha_c}{2} [\partial_\mu \partial_\nu h_{\rho\sigma} + \partial_\rho \partial_\sigma h_{\mu\nu} - (\mu \leftrightarrow \nu)]$ and to the order $\alpha_c^2$ results in

$$\sqrt{-g} = 1 + \frac{\alpha_c}{2} h_{\mu\nu}^2 + \frac{\alpha_c^2}{8} (h_{\mu\nu}^2 h_{\rho\sigma}^2 - 2h_{\mu\rho}^2 h_{\sigma\nu}^2) ,$$

$$\sqrt{-g} R_{\mu\nu,k} = \frac{\alpha_c^2}{4} (\partial_\mu \partial_\nu h^\lambda + \partial_\lambda \partial_\mu h^\nu - \partial_\nu \partial_\mu h - \partial^2 h_{\mu\nu})^2 ,$$

$$\sqrt{-g} R^2 = \alpha_c^2 (\partial^2 h - \partial_\mu \partial_\nu h_{\mu\nu})^2 . \tag{32}$$

With this the weak-field expansion of the Weyl action $(\mu\nu\alpha\beta)$ reads (modulo total derivatives) \[21\]

$$\mathcal{A}_{\text{conf}} = -\frac{1}{8\alpha_c^2} \int d^4x (g)^{1/2} C_{\mu\nu\kappa\lambda} \nabla_{\mu\nu\kappa\lambda} = -\frac{1}{16} \int d^4x \, h^{\alpha\beta} P^{(2)}_{\mu\nu,\alpha\beta} \delta^2 h_{\alpha\beta} = -\frac{1}{16} \int d^4x (\nabla^2 h^{\alpha\beta})^2 . \tag{33}$$

Here, $P^{(2)}_{\mu\nu,\alpha\beta} = \pi_{\mu\nu}(\pi_{\beta}) - \frac{3}{2} \pi_{\mu\nu}(\pi_{\alpha\beta})$ is the spin-2 projection, and $\pi_{\mu\nu} = \eta_{\mu\nu} - \partial_\mu \partial_\nu$ is the transverse vector projection. We also used $h^\alpha\beta = h_{\mu\nu}^\alpha + \frac{1}{2} \eta_{\mu\nu} \varphi$ ($h^\alpha\beta = 0$ so that $h_{\alpha\beta} = 0$) and b) $h_{\mu\nu} = \hbar_{\mu\nu} + \partial_\mu \eta_{\nu} - \partial_\nu \eta_{\mu} + \partial_\mu \partial_\nu - \frac{3}{2} \eta_{\mu\nu} \partial^2 \sigma$ (with $\partial_\mu h^{\alpha\beta}_\mu = 0$ and $\partial^2 \eta_{\mu\nu} = 0$) which serve to identify irreducible degrees of freedom. Using the conformal gauge $\varphi = 0$ and the coordinate gauge $\partial_\mu h^{\mu\nu} = 0$ (with the associated Faddeev–Popov operator $(M_{\text{FP}})_{\mu\nu} = -\nabla_\mu \eta_{\nu} - \eta_{\mu\nu}$) the functional measure reads

$$Dh_{\mu\nu} = D\hbar_{\mu\nu} D\eta_{\mu\nu} D\sigma D\varphi \det(-\nabla_\mu \eta_{\nu})^{1/2} \det(-\eta_{\mu\nu})_{\eta_{\mu\nu}}^{1/2} \rightarrow D\hbar_{\mu\nu} \det(-\nabla_\mu \eta_{\nu})^{1/2} \det(-\eta_{\mu\nu})_{\eta_{\mu\nu}}^{1/2} . \tag{34}$$

For one-loop effective action calculations we use the more customary approach in which the weak-field action $\mathcal{A}_{\text{conf}}$ is written in terms of unconstrained variable $h_{\nu\lambda}$ as

$$-\frac{1}{16} \int d^4x \left[ h^{\mu\nu} \nabla^2 h_{\mu\nu} - \partial_\mu \lambda h^\lambda h_{\mu\nu} \partial_\nu h_{\mu\nu} - \frac{1}{6} \hbar^2 \right] \tag{35}$$

Here $H_{\mu\nu} = 1/2 \partial_\mu \partial_\nu - \eta_{\mu\nu}$ and $h = h_{\mu\nu} - \partial_\mu \partial_\nu h_{\mu\nu}$. To obtain the diagonal kinetic operator one has to cancel the second and third term by fixing the gauges: $\lambda = \partial_\mu h^{\mu\nu} = \zeta'(x)$ (coordinate gauge) and $\chi \equiv \hbar = \chi(x)$ (conformal gauge). The Faddeev–Popov operator for coordinate gauge is $(M_{\text{FP}})_{\mu\nu} = -\nabla_\mu \eta_{\nu} - \eta_{\mu\nu}$ while for conformal gauge $(\Lambda_{\text{FP}}) = (D - 1) \delta^{(D)}(x - y)$. In this case the functional-integral measure is:

$$Dh_{\mu\nu} \rightarrow Dh_{\mu\nu} \delta[\chi - \zeta][\chi' - \zeta'] \det(M_{\text{FP}}) . \tag{36}$$

($\zeta$ and $\zeta'$ are arbitrary functions of $x$). With the help of ’t Hooft’s averaging trick the corresponding partition function coincides with that obtained from \[33\] and \[34\].
8

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