Ultraviolet Completion of the Big Bang in Quadratic Gravity

Ruolin Liu, [1,2,3,*] Jerome Quintin, [4,5,6,2,3,†] and Niayesh Afshordi [1,2,3,‡]

[1] *Department of Physics and Astronomy, University of Waterloo, Waterloo, Ontario N2L 3G1, Canada*
[2] *Perimeter Institute for Theoretical Physics, Waterloo, Ontario N2L 2Y5, Canada*
[3] *Waterloo Centre for Astrophysics, University of Waterloo, Waterloo, Ontario N2L 3G1, Canada*
[4] *École de technologie supérieure, Université du Québec, Montréal, Quebec H3C 1K3, Canada*
[5] *Department of Physics, McGill University, Montréal, Quebec H3A 2T8, Canada*
[6] *Department of Applied Mathematics, University of Waterloo, Waterloo, Ontario N2L 3G1, Canada*

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We present a quantum quadratic gravity inflationary scenario that can accommodate the new cosmological constraints, which have disfavored Starobinsky inflation. The theory is asymptotically free in the ultraviolet, but 1-loop running is found to dynamically lead to slow-roll inflation toward the infrared. When a large number of matter fields contribute to the beta functions, the spectral index and the tensor-to-scalar ratio can be phenomenologically viable. We find that as inflation ends, the theory approaches its strong coupling regime and general relativity must emerge, as an effective field theory, as the universe must reheat and enter its standard radiation era. In order to avoid strong coupling, a minimum tensor-to-scalar ratio of 0.01 is predicted for this theory. Our framework offers a laboratory for connecting a concrete ultraviolet completion (quantum quadratic gravity) with inflationary dynamics, reheating, and precise cosmological observations.

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General relativity (GR) is tremendously successful when treated as an effective field theory (EFT) [1]. This means, though, that there is a cutoff scale beyond which it cannot be trusted. Beyond this scale, one has to deal with GR's nonrenormalizability and with singularities that develop in solutions to the classical Einstein equations. When starting with the linear-in-curvature Einstein-Hilbert (EH) action [2], $S_{\mathrm{EH}} = \frac{1}{2}\int \mathrm{d}^4x\sqrt{-g}M_{\mathrm{Pl}}^2 R$, one has to include an infinite hierarchy of higher curvature counterterms to cancel divergences in large-momentum loops [3–5], but this is not the case if the bare action is augmented only by terms that are quadratic in curvature [6]. As such, including quadratic gravity,

$$S_{\text{quadratic gravity}} = -\int \mathrm{d}^4x\sqrt{-g}\left(\frac{R^2}{\xi} + \frac{C^2}{2\lambda}\right), \qquad (1)$$

allows for a possible ultraviolet (UV) completion of GR. It is thus compelling to explore UV regimes where the effects of quadratic gravity could manifest themselves and be tested, such as in black holes [7,8] and in the early universe (e.g., [9]). This Letter explores the latter.

When adding the square of the Ricci scalar (R^2) to the EH action, one already has the starting point of Starobinsky inflation [10]. However, in general, quadratic gravity also contains the contraction of the Weyl tensor with itself (C^2) [11]. Hence, as a higher-derivative theory, the theory comes with two new fields [9,12] in addition to the massless graviton in GR: a massive spin-0 field coming from R^2 (which can act as the inflaton) and a massive spin-2 field coming from C^2. The latter field is a ghost; i.e., it carries negative kinetic energy, leading to an unbounded Hamiltonian (if we ignore interactions). There is a lot of literature attempting to make sense of ghosts in various settings [9,13–32], which we do not review here, but in the context of inflationary cosmology different techniques allow one to handle the ghost, its quantization, and the corresponding analysis of Starobinsky-like inflation, where the presence of C^2 only leads to (small) corrections to the tensor perturbations [33–40].

Instead, here we wish to explore another possibility—arguably more radical—distinct from Starobinsky's $R + R^2$ inflation. As quantum quadratic gravity (QQG) is UV complete, one may start at the "big bang" singularity at infinite curvature. There, QQG is "pure"—there is no GR. In this regime, the theory is amenable to a perturbative treatment, and accordingly, beta functions at the 1-loop order reveal asymptotic freedom in the UV [41–46]. QQG

*Contact author: r388liu@uwaterloo.ca
†Contact author: jquintin@physics.mcgill.ca
‡Contact author: nafshordi@pitp.ca

has many similarities to quantum chromodynamics (QCD) and suggests that GR may emerge as QQG becomes strongly coupled in the infrared (IR) [47,48] (see also the recent proposal [49]); from the low energy perspective, QQG would be confined. The question thus becomes, can one find a successful early universe scenario in the deep UV, solely from QQG, i.e., quadratic gravity with running taken into account?

We should highlight that what would make this scenario unique is that gravity does become strongly coupled at some scale, which is also used to confine ghost degrees of freedom [47,48]. In contrast, in Starobinsky inflation, gravity remains weakly coupled, and thus the containment of ghosts and Ostrogradsky instabilities would require a more creative solution.

The renormalization group (RG) flow of QQG and other higher-derivative theories has been the subject of renewed interest in recent years [50–54]. In particular, [50] argued that the previously derived 1-loop beta functions of QQG did not capture the physical scale at which interactions take place. New beta functions have been computed,

$$\beta_\xi = \frac{\mathrm{d}\xi}{\mathrm{d}\ln\mu} = -\frac{1}{(4\pi)^2}\frac{\xi^2 - 36\lambda\xi - 2520\lambda^2}{36},$$
$$\beta_\lambda = \frac{\mathrm{d}\lambda}{\mathrm{d}\ln\mu} = -\frac{1}{(4\pi)^2}\frac{[(1617+90\mathcal{N})\lambda - 20\xi]\lambda}{90}, \quad (2)$$

where μ represents here the *physical* running scale of the QQG couplings ξ and λ. Although the validity of these "physical" beta functions is still a matter of active debate [51–54], in this Letter we shall argue that this formulation [in particular, the emergence of a maximum for $\xi(\mu)$] would allow for a successful inflationary scenario. Thus, we use this observation as empirical guidance that may inform the ongoing theoretical debate.

While [50] derived these equations in vacuum, we added the possible contribution from matter fields in the loops; the number of such fields,

$$\mathcal{N} = \frac{1}{60}\mathcal{N}_{\text{scalar}} + \frac{1}{5}\mathcal{N}_{\text{vector}} + \frac{1}{20}\mathcal{N}_{\text{fermion}}, \quad (3)$$

correspondingly enhances β_λ [9]. Exact analytic (though implicit) solutions can be derived in the vacuum ($\mathcal{N} = 0$) case [55]. The general features remain the same even as $\mathcal{N}$ is increased; we show the case of $\mathcal{N} = 10$ in Fig. 1 for illustrative purposes. The key aspect of the results of [50] is that the beta functions now admit solutions that are both asymptotically free in the UV (reaching the green star at the origin in Fig. 1) and tachyon free in the IR, which is to say that low-μ basin of attraction is where $\lambda > 0$ and $\xi < 0$ (the pale blue region in Fig. 1); the orange curve is an example of such a trajectory. Starobinsky inflation could happen precisely in the tachyon-free region, should one include an EH term in the action (see Ref. [56]).

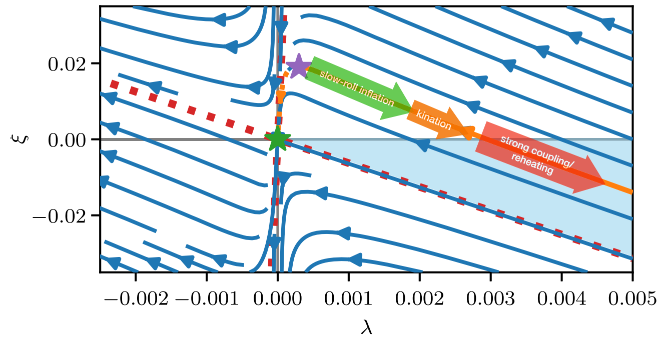

FIG. 1. The renormalization group (RG) flow of quadratic gravity with $\mathcal{N} = 10$ matter-field content. The two red dashed lines show the separatrices of the RG flow, while the light blue shaded region is tachyon-free [50]. The orange trajectory shows a candidate cosmological evolution: starting from a potential tunneling from a (no-boundary) Euclidean geometry at the top of the trajectory (purple star) → slow-roll inflation (potential-dominated) → kination (kinetic-dominated) → strong-coupling and reheating, where an EH term with effective M_{Pl} (i.e., GR) emerges eventually.

The scenario we wish to put forward is different, as we wish to start with pure QQG in the UV. Looking at trajectories that would eventually reach the tachyon-free region (where we would expect GR to emerge), we see that all of them (except the one trajectory that follows the red dotted separatrix) must first go in the region where both λ and ξ are positive: λ always monotonically grows, while ξ first grows before decreasing and eventually crossing the tachyon divide at a scale that we call μ_0. What we find is that a successful inflationary phase can be realized in the regime where ξ is decreasing, though before it crosses the tachyon divide. We find that as inflation ends, the universe enters a kinetic-dominated phase (kination) and QQG approaches its strong coupling scale. Eventually, GR must emerge, and the universe must enter its radiation era. One can thus think of the reheating surface at the end of kination following inflation as a UV-IR matching surface for QQG-GR.

Here is how inflation arises in pure QQG: considering a homogeneous and isotropic background, there is no contribution from the Weyl tensor, so the action is solely that of R^2 gravity. Without running, the theory is scale invariant [57,58], and any constant-R spacetime extremizes the action, such as Minkowski and (anti-)de Sitter. The exact scale invariance is broken once a physical scale appears in the theory, such as M_{Pl} if one includes GR; that is how one goes from de Sitter to quasi–de Sitter in Starobinsky inflation. However, even without such explicit scale, something similar occurs if we consider the running due to quantum effects: the theory is not exactly pure R^2 gravity anymore, and scale invariance is broken, akin to the so-called quantum *conformal anomaly* [59,60].

While the onset of inflation in this scenario (like all the alternatives) remains speculative, a natural possibility is a no-boundary Euclidean manifold [61,62]. Such a manifold would be an exact solution to the Euclidean QQG at the maximum of ξ (where its derivative vanishes, depicted by the purple star in Fig. 1), as the Euclidean action for a 4-sphere is $S_{\rm E} \propto \xi^{-1}$ in this theory. Across its equator, the Euclidean half-sphere can then be matched to the waist of a closed (Lorentzian) de Sitter spacetime, which would describe cosmology when $t \to -\infty$. However, this de Sitter phase is unstable due to RG running of ξ, which is expected to initiate slow-roll inflation [63].

To explore the effects of RG flow, we start with one important assumption, which is that we are not in vacuum, but that instead a large number of matter fields are present. This is usually a fair assumption in most models that go beyond the standard model of particle physics [64]. This is not to say that such matter fields are necessarily excited at the big bang; in fact, we assume that they sit in their vacua. Still, vacuum fluctuations of $\mathcal{N} \gg 1$ fields have important implications in the running of QQG, and thus indirectly, these matter fields contribute to modifying the cosmological background and perturbations. The solution to the beta functions (2) in the large-$\mathcal{N}$ limit can be expressed as [55]

$$\xi(\mu) \simeq \frac{35\lambda_0^2 \ln(\mu/\mu_0)}{8\pi^2[1 + \lambda_{\rm tH}\ln(\mu/\mu_0)]}, \tag{4}$$

where μ_0 is the scale where $\xi = 0$ and $\lambda = \lambda_0$ [65]. Here, we define a 't Hooft–like coupling constant, $\lambda_{\rm tH} \equiv \lambda_0 \mathcal{N}/(4\pi)^2$, which is the key quantity in assessing the strength of the loop corrections when $\mathcal{N} \gg 1$; the 1-loop approximation is under control so long as $\lambda \lesssim \lambda_0$ and $\lambda_{\rm tH} \lesssim 1$, so λ_0 must be sufficiently small at large $\mathcal{N}$. We shall see that as λ flows beyond λ_0 (as we cross the tachyon divide), strong coupling must be reached and GR must emerge as $\lambda\mathcal{N}/(4\pi)^2$ surpasses unity.

At this point, different techniques can be employed to compute the quantum effective action [67] that results from the running of the couplings ξ and λ. We take the simple approach of promoting μ to a covariant definition of energy scale and substituting the resulting coupling into the action. This approach is ambiguous, though, to the extent that it depends on the choice of energy scale for μ, but there are many examples (including in inflation; e.g., [68–71]) where it is reasonable to expect the RG scale to correspond to the background cosmological or curvature scale. As such, many curvature invariants could be chosen, e.g., the Kretschmann scalar, but we take what is possibly the simplest: the Ricci scalar [72]. Considering the action of QQG without the Weyl tensor first [73], this means that we can write the action as an $f(R)$ theory of gravity; taking the solution (4), it is $f(R) = R^2/\xi(R) = 8\pi^2(\lambda_{\rm tH} + 4/\ln[R^2/\mu_0^4])R^2/(35\lambda_0^2)$. The RG flow thus yields a small (logarithmic) correction to the pure R^2 action, which leads to quasi–de Sitter attractor solutions [55].

As observables at this level should not depend on the choice of frame (e.g., [36]), let us transform to the Einstein frame, where more intuition can be gained since the $f(R)$ theory becomes equivalent to GR with a scalar field. Sufficiently far from the end of inflation, the scalar field's potential can be approximated as [55]

$$V(\varphi) \simeq \frac{35\lambda_0^2\mu_0^4}{128\pi^2\lambda_{\rm tH}}\left(1 - \frac{\sqrt{6}\mu_0}{\lambda_{\rm tH}\varphi}\right). \tag{5}$$

This falls in the brane inflation generalized *phenomenological* classification of [74], but as far as we are aware, it is the first time that this specific form of the potential is derived from a UV-complete theory. The potential approximated as (5) only holds at early times, but it nevertheless still captures an important finding about late times [55], which is that inflation has a graceful exit mechanism as the potential steepens, but also that it does not reach a minimum where the inflaton could decay. Instead, the field rolls ever faster into a kinetic-dominated phase (kination). Eventually, though, nonperturbative effects would alter the dynamics as the theory enters its strong coupling regime.

Assuming that the number of e-folds N in between the end of inflation and the time of horizon exit for modes of cosmological interest is sufficiently large (50 to 60, similar to other inflationary models [75]; see Ref. [55] for an assessment of this assumption), we find the amplitude of the scalar perturbations to be given by $A_{\rm s} \sim 35\lambda_0^2 N^{4/3}/[512\pi^4(2\lambda_{\rm tH})^{1/3}]$, while the scalar spectral index and the tensor-to-scalar ratio are, respectively, [55]

$$n_{\rm s} \sim 1 - \frac{4}{3N}, \qquad r \sim \frac{8}{3}\left(\frac{2}{\lambda_{\rm tH}^2 N^4}\right)^{1/3}. \tag{6}$$

This somewhat resembles the predictions of Starobinsky inflation, though both are larger than Starobinsky's $n_{\rm s} \approx 1\text{–}2/N$ and $r \approx 12/N^2$.

Using the exact expressions derived in [55], we can plot the predictions in the $n_{\rm s}$-r plane (see Fig. 2), where this is also contrasted with Starobinsky inflation. Dots and lines of different color correspond to different values of the 't Hooft–like coupling $\lambda_{\rm tH}$ in QQG. We can see that $\lambda_{\rm tH}$ has little effect on $n_{\rm s}$, but the more weakly coupled the theory is (so the smaller $\lambda_{\rm tH}$ is), the larger r. In comparison, we show a combination of constraints from cosmic microwave background (CMB) data (from *Planck* [76], ACT [77], SPT [78], and BICEP+Keck [79][BK]) and baryon acoustic oscillation (BAO) data (from DESI [80]). Such data combination (especially ACT with DESI within ΛCDM; see Ref. [81] for a discussion of the tension) prefers larger $n_{\rm s}$ values, putting Starobinsky inflation in slight tension with observations but placing our QQG model in a favorable position. Nonetheless, allowing for a dynamical dark energy [82] would put both Starobinsky inflation and QQG within 1σ of observations [78].

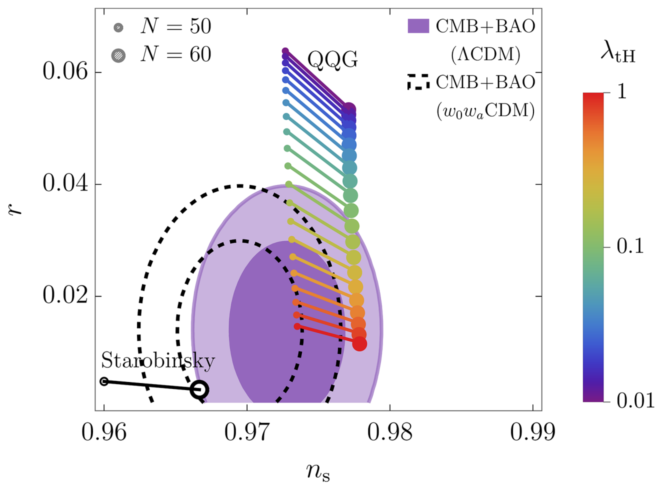

FIG. 2. Plot of the scalar spectral index $n_{\rm s}$ vs the tensor-to-scalar ratio r for Starobinsky inflation (in black), contrasted with our QQG model in color for different values of the coupling $\lambda_{\rm tH}$. The purple and dashed contours correspond to a combination of CMB (*Planck*18 + ACT + SPT + Lensing + BK) and BAO (DESI) observational constraints for ΛCDM and w_0w_aCDM cosmologies, respectively [78].

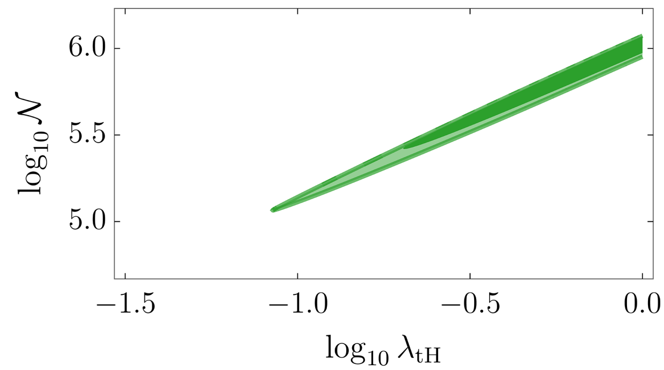

FIG. 3. Constraints on the QQG parameters $\lambda_{\rm tH}$ and $\mathcal{N}$ from the observationals bounds on $n_{\rm s}$, r, and $A_{\rm s}$ combined. The dark and pale green regions amount to 1σ and 2σ confidence intervals, respectively.

Given a fixed number of inflationary e-folds, the observables $n_{\rm s}$, r, and $A_{\rm s}$ solely depend on N, λ_0, and $\lambda_{\rm tH} \equiv \lambda_0\mathcal{N}/(4\pi)^2$. Thus, observational constraints on those quantities can be translated into bounds on $\lambda_{\rm tH}$ and $\mathcal{N}$. We present such constraints in Fig. 3, where we set the scalar amplitude to $A_{\rm s} = e^{3.06} \times 10^{-10}$ [76], and 1σ and 2σ confidence intervals on $n_{\rm s}$ and r are taken as in Fig. 2, while for the e-folding number we take 50 ± 10. Since $r \sim \lambda_{\rm tH}^{-2/3}$, we can see in Fig. 3 that there is no upper bound on $\lambda_{\rm tH}$ because there is no observational lower bound on r. However, we do not trust the results beyond $\lambda_{\rm tH} \gtrsim 1$ (in the strong coupling regime), hence the cutoff at $\lambda_{\rm tH} = 1$. As such, the viable parameter space is spanned by $0.1 \lesssim \lambda_{\rm tH} \lesssim 1$ and $10^5 \lesssim \mathcal{N} \lesssim 10^6$. It thus appears that two things are phenomenologically preferred: that there must be a (very) large number of matter fields present in the theory; and that the theory must be very close to strong coupling as λ approaches the tachyon divide at λ_0 since $\lambda_{\rm tH} \equiv \lambda_0\mathcal{N}/(4\pi)^2$ is likely very close to unity. The latter brings us back full circle to our earlier claim: current observations suggest that crossing the tachyon divide appears to be nearly coincident with entering the strong-coupling regime, $\lambda_{\rm tH} \gtrsim 1$, as well as with the onset of reheating (red arrow in Fig. 1).

Let us summarize our findings. We presented a UV-complete “quantum quadratic gravity” inflationary scenario that can be compatible with the recent CMB constraints, which may otherwise disfavor the standard Starobinsky $R + R^2$ inflation. By identifying the physical renormalization scale with curvature, $\mu = |R|^{1/2}$, pure R^2 gravity acquires controlled logarithmic corrections and yields an almost-plateau potential in the Einstein frame. This quantum quadratic gravity preserves the slow-roll virtues of plateau models while shifting predictions toward the parameter space preferred by recent CMB observations. In particular, for a large number of matter fields, $\mathcal{N} \sim \mathcal{O}(10^5\text{–}10^6)$, the ’t Hooft–like coupling $\lambda_{\rm tH}$ can stay under perturbative control, and the predicted $\{n_{\rm s}, r, A_{\rm s}\}$ are consistent with the data. In order to avoid strong coupling, the tensor-to-scalar ratio r is predicted to be $\gtrsim 0.01$.

Many things remain to be explored. First, incorporating two-loop RG equations for the quadratic couplings would test the robustness of the quantum inflation potential (especially as we approach $\lambda_{\rm tH} \sim 1$), refine the viable parameter space, and quantify scheme and threshold uncertainties as GR effectively emerges. The RG flow of other couplings, such as Gauss-Bonnet, Yang-Mills, and Yukawa, may lead to richer phenomenology as observations improve (e.g., [85]). Second, while the Weyl term C^2 vanishes on homogeneous and isotropic backgrounds, it affects perturbations; assessing its impact—together with possible ghost mechanisms and the ensuing stability and observables—is an important next step. One can also think of comparing this RG-improved setup with holographic cosmology (which also requires loop contributions with a large value of $\mathcal{N}$ [85,86]) and the Hartle-Hawking no-boundary proposal, clarifying the role of initial state in our QQG inflation. Moreover, a kinetic-dominated (kination) phase naturally appears at the end of inflation; analyzing the onset of reheating, the strong-coupling window, and the EFT handover to GR could sharpen the link to data. Finally, beyond ACT, a combined analysis with *Planck*+BICEP+SPT+Simons Observatory could probe the reported ACT tension [81] and forecast r within next-generation sensitivities. Overall, this framework offers a concrete laboratory connecting RG running in renormalizable quadratic gravity to inflationary dynamics, reheating, and precision CMB data, while inviting broader comparisons with UV-complete cosmological scenarios.

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Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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[64] One may question how large $\mathcal{N}$ can reasonably be. In theoretical settings, such as holography, $\mathcal{N}$ is often taken to be arbitrarily large, to ensure perturbative gravity in the bulk. Here, we remain agnostic about whether or not a very large $\mathcal{N}$ represents a fine-tuning problem. We should further note that these matter fields would probably end up being confined to the UV, together with all the degrees of freedom of QQG, in a similar fashion to quarks and gluons in QCD.

[65] Here, we ignore the impact of the running of the coupling to Gauss-Bonnet's invariant $\mathcal{G}$, which is not widely studied due to its topological nature. However, this may affect cosmological evolution—as in $f(R,\mathcal{G})$ theories, which are widely studied in the literature, e.g., [66]—but we defer that consideration to future work.

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[73] The contribution from the Weyl tensor vanishes on a homogeneous and isotropic background. This is not true once cosmological perturbations are introduced. Still, contributions from the Weyl tensor to, e.g., tensor perturbations are usually highly suppressed [34], and we do not expect this result to be affected by small logarithmic running.

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