

Gravitational Origin of the QCD Axion

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Gravity can give rise to (pseudo)scalar fields—for instance due to torsion. In particular, axions of gravitational origin have been proposed as a minimal and compelling solution to the strong CP problem. In this Letter, we critically examine the feasibility of this proposal. We demonstrate that models in which the scalar field couples to fermionic currents only through derivatives do not yield a satisfactory axion. Moreover, we identify the necessary conditions for generating a gravitational axion through quantum effects, highlighting Weyl-invariant Einstein-Cartan gravity as a promising theoretical setting.

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The strong CP problem—At present, all known phenomena at microscopic scales can be accurately described by the standard model of particle physics (SM). Having withstood extensive experimental scrutiny, the SM was formally completed with the discovery of the Higgs boson [1,2]. Arguably, however, our understanding of the SM is still far from complete. Among the most important open questions is the strong CP problem: the puzzling experimental fact that quantum chromodynamics (QCD), the theory of strong interactions, preserves CP symmetry to an extraordinary accuracy despite having no apparent reason to do so.

A compelling resolution to the strong CP problem relies on the introduction of a new pseudoscalar particle—the axion [3–5]—that dynamically suppresses CP violation in QCD (see reviews [6,7]). In addition to solving the strong CP problem, the axion is a well-motivated dark matter candidate and thereby addresses a crucial open issue in cosmology. Extensive experimental efforts are underway to detect the axion [8–17]; for an overview, see Ref. [18].

Axion origins—In the oldest proposals, the axion arises from the spontaneous breaking of a global Peccei-Quinn symmetry [3–5]. However, the most popular models [19–22] require the introduction of at least 5 degrees of freedom in addition to the axion. Moreover, global symmetries are generally expected to be violated by

quantum gravitational effects (see, e.g., Ref. [23]), posing a challenge to the theoretical consistency of this approach, as Planck-suppressed operators could reintroduce CP violation [24–26].

These concerns have motivated the development of alternative axion models, including those based on extra dimensions [27] (see review [28]), or constructions where the axion is the Stückelberg field of the QCD gauge redundancy [29–32]. Rather than being an obstacle, gravity itself may play a constructive role: the axion could be of purely gravitational origin [33–39]. This scenario is especially appealing due to its minimality, requiring no ingredients beyond the SM and general relativity (GR). Furthermore, recent work indicates that the existence of an axion is a requirement for the consistency of quantum gravity itself [32,40,41].

In this Letter, we critically examine the proposal that the axion originates from purely gravitational dynamics. In the first part, we show that previous arguments—as used in [33–39]—are not successful. Our proof is almost trivial and relies on the fact that the relevant operator [shown in Eq. (7)], as considered in existing proposals, fails to break the shift symmetry and hence does not generate any potential. If a scalar obeys an exact shift symmetry, it is unable to dynamically solve the strong CP problem and hence the corresponding particle cannot be an axion.

In a second, admittedly more speculative part, we develop approaches for saving the axion of gravitational origin. We demonstrate that, generically, the “problematic” shift symmetry is broken in Einstein-Cartan (EC) gravity. Nevertheless, the (perturbatively generated) interactions between the axion and quarks and gluons turn out either unrealistically weak or hinge on severe tunings; therefore, the strong CP problem persists. Finally, we point out that the Weyl-invariant subclass [39] of EC gravity is special. The requirement of exact Weyl symmetry at the

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quantum level allows the control of radiative corrections and leads to a tower of higher-dimensional operators that can enable the gravitational scalar to assume the role of the QCD axion. Thus, even though the original argumentation of [39] was not accurate, the theory of [39] may lead to a successful solution of the strong CP problem.

Review of QCD axion—At the heart of the strong CP problem lies the operator

$$\mathcal{L} \supset \frac{\alpha}{4\pi} \bar{\theta} \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}, \quad (1)$$

where $\alpha = g^2/(4\pi)$ with g the QCD gauge coupling and $\bar{\theta}$ is the physical CP violating parameter, i.e., $\bar{\theta}$ is the sum of the vacuum superselection angle of QCD and the CP violating phase of the quark mass matrix. Moreover, $G_{\mu\nu}$ and $\tilde{G}_{\mu\nu}$ are the field strength tensor of QCD and its dual, respectively, and the trace is taken over color indices.

Although Eq. (1) is a total derivative and hence does not contribute perturbatively to equations of motion, it causes nonperturbative CP violation [20,42,43]. In turn, this translates in the neutron having an electric dipole moment $d_n \simeq 10^{-16} \bar{\theta}$ e cm; when compared with the experimentally measured value [44] $d_n^{\text{exp}} \lesssim 10^{-26}$ e cm, it dictates that

$$\bar{\theta} \lesssim 10^{-10}. \quad (2)$$

The strong CP problem corresponds to the question why $\bar{\theta}$ is so small.

The key of the axion solution [3–5] is to effectively promote $\bar{\theta}$ to a dynamical variable—a canonically normalized pseudoscalar field a —that extends Eq. (1) as

$$\mathcal{L} \supset \frac{\alpha}{4\pi} \left(\bar{\theta} - \frac{a}{f_a} \right) \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}, \quad (3)$$

where f_a is the axionic decay constant. Since $\text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}$ is a total derivative, the coupling (3) leads to a perturbative shift symmetry for the axion. Nonperturbative QCD effects break this shift symmetry and induce a potential for the axion [3–5], which can be well approximated as [45,46] (see also discussion in [7])

$$V = -m_\pi^2 f_\pi^2 \sqrt{1 - \frac{4m_u m_d}{(m_u + m_d)^2} \sin^2 \left[\frac{1}{2} \left(\bar{\theta} - \frac{a}{f_a} \right) \right]}. \quad (4)$$

Here, m_π and f_π are the pion mass and decay constant, respectively, while m_u and m_d correspond to the masses of the up and down quarks. Evidently, Eq. (4) is minimized for $a/f_a = \bar{\theta}$ [47], and so the axion resolves the strong CP problem. Key for achieving this is the breaking of the axionic shift symmetry by QCD effects.

It should be emphasized that we identify as QCD axion any particle that receives its main contribution to the

potential from the operator (3), irrespectively of its UV completion.

Review of gravitational axion—Although the concrete implementations [33–39] of how to obtain an axion from gravity [33–36] differ, the essence of the idea is the following. The starting point is the fact that GR exists in different formulations (see Refs. [48–50] for overviews). All of them are fully equivalent in the absence of matter and for a simple choice of gravitational action, e.g., consisting only of the Einstein-Hilbert term.

Most commonly used is the metric—or second-order—formulation, in which spacetime is fully described by curvature. However, strong arguments speak in favor of the EC formulation, which is endowed with torsion in addition to curvature [51–54]. In particular, EC gravity can be derived from gauging the global Lorentz group [55–57], which puts the gravitational interaction on the same footing as the other SM forces. Related to the gauging is the fact that coupling fermions to GR naturally leads to the EC formulation of GR [56] (see also Refs. [58,59]).

In EC gravity, the minimal kinetic term of a spinor Ψ splits as (see Refs. [58,60])

$$\frac{i}{2} \bar{\Psi} \gamma^\mu \mathcal{D}_\mu \Psi + \text{H.c.} = \left(\frac{i}{2} \bar{\Psi} \gamma^\mu \mathring{\mathcal{D}}_\mu \Psi + \text{H.c.} \right) - \frac{1}{8} J_A^\mu \hat{T}_\mu, \quad (5)$$

where $\mathcal{D}_\mu$ is the fermionic covariant derivative, $\mathring{\mathcal{D}}_\mu$ denotes its torsion-free, i.e., Riemannian, counterpart, and we introduced the axial fermionic current

$$J_A^\mu = \bar{\Psi} \gamma^5 \gamma^\mu \Psi. \quad (6)$$

Moreover, γ^μ and γ^5 are gamma matrices and $\hat{T}_\alpha = \epsilon_{\alpha\beta\mu\nu} T^{\beta\mu\nu}$ corresponds to the axial torsion vector, with $T_{\alpha\nu}^\mu$ derived from the connection $\Gamma_{\alpha\nu}^\mu$ as $T_{\alpha\nu}^\mu \equiv 1/2(\Gamma_{\alpha\nu}^\mu - \Gamma_{\nu\alpha}^\mu)$; see Appendix A of Supplemental Material [61] for more details about conventions.

As is well-known, cf., e.g., [39,59,62–70], axial torsion contains a pseudoscalar degree of freedom ϕ , which can be heuristically extracted as $\hat{T}_\alpha = \partial_\alpha \phi$. As evident from Eq. (5), ϕ couples to all fermions via

$$\mathcal{L} \supset J_A^\mu \partial_\mu \phi. \quad (7)$$

Now we can integrate by parts and use the nonconservation of the chiral current, $\mathring{\mathcal{D}}_\mu J_A^\mu \supset \alpha/4\pi \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}$, to conclude that the action contains a term

$$\mathcal{L} \supset \frac{\alpha}{4\pi} \phi \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}. \quad (8)$$

When this is compared with Eq. (3), it appears that ϕ couples to the topological density $\text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}$ in exactly the same way as the QCD axion (with the identification $\phi = a/f_a$). Therefore, one is immediately tempted to

conclude that ϕ can assume the role of an axion, as was done in [33–39].

For completeness, we mention that this approach can only succeed if ϕ corresponds to a propagating degree of freedom, i.e., if it possesses a kinetic term; note that EC gravity by itself does not contain dynamical torsion, and thus it contains no other modes besides the massless graviton. A self-consistent way to enrich the particle spectrum is to supplement the action with invariants that are quadratic in curvature [71]. For the case at hand, the appropriate term to be added is $\tilde{R}^2$ [39], where $\tilde{R}$ is the pseudoscalar (Holst) curvature defined in Appendix B of Supplemental Material [61]. Since this point is inessential for our investigation, we shall not discuss it further and refer the interested reader to the aforementioned article and references therein.

Proof to the contrary—We will now show why the seemingly innocuous argument above does not go through. To this end, it suffices to realize that before any manipulations, the operator $\partial_\mu \phi J_A^\mu$, as shown in Eq. (7), is *manifestly* and *exactly* shift-symmetric. No effect, perturbative or nonperturbative, can break this symmetry, unless the correlator of J_A^μ has a pole at zero momenta, and this is not the case here.

Thus, the coupling (7) does not lead to a potential for ϕ [82], and without a potential, there is no mechanism to dynamically relax the CP violating angle to zero. Hence, the operator (7) does not suffice to turn ϕ into an axion.

What has gone wrong?—What remains to be done is demonstrate that the argument based on Eq. (8) is incomplete. The key point is that there is another contribution to the chiral anomaly from the mass m_Ψ of the fermion (see, e.g., Refs. [7,84]),

$$\hat{D}_\mu J_A^\mu = \frac{\alpha}{4\pi} \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu} + 2i\bar{\Psi} m_\Psi \gamma^5 \Psi. \quad (9)$$

Therefore, the term $\phi \hat{D}_\mu J_A^\mu$ entails two contributions to the axionic potential: one from the interaction with gluons and the other from the Yukawa coupling to the fermions. Consistency with the exact shift symmetry demands that they cancel exactly [85].

This argument can also be made by considering the (anomalous) axion-dependent chiral rotation

$$\Psi_{L,R} \mapsto e^{\mp i\phi/2} \Psi_{L,R}, \quad (10)$$

where $\Psi_{L,R} = 1/2(1 \mp \gamma_5)\Psi$. In this way, one trades the derivative coupling $\partial_\mu \phi J_A^\mu$ for the following new terms:

$$\mathcal{L} \supset -\left(e^{-i\phi} \bar{\Psi}_L m_\Psi \Psi_R + \text{H.c.}\right) + \frac{\alpha}{4\pi} \phi \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu}, \quad (11)$$

which, when expanded around $\phi \approx 0$, give rise to exactly the same contributions to the two-point function of ϕ as Eq. (9).

A few remarks are in order at this point. First, it should not come as a surprise that the fermion mass plays a crucial role. Indeed, the topological term (1) does not lead to observable CP violation if at least one massless fermion is present [3,42,43,86,87]. Second, one can wonder if an analogous argument can be made for the “true” axion coupling in Eq. (3). Since $\text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu} = \partial_\mu K^\mu$ with Chern-Simons current K^μ , one could try to integrate by parts and rewrite

$$a \text{Tr} G_{\mu\nu} \tilde{G}^{\mu\nu} \propto -K^\mu \partial_\mu a. \quad (12)$$

So one might be tempted to conclude that, also here, the axion obeys an exact shift symmetry. However, the crucial difference with Eq. (7) is that K^μ is not gauge-invariant, and so it is impossible to define the fundamental Lagrangian respecting the local $SU(3)$ symmetry in terms of the Chern-Simons current. In addition, the correlator of K^μ has a (gauge-invariant) pole at zero momenta [88].

Quantum corrections—We have given a rigorous proof showing that previous arguments for generating an axion of gravitational origin—as used in [33–39]—are not successful. Now we come to the second, more speculative, part of the Letter, in which we develop approaches for saving the gravitational axion.

In fact, there is a loophole in the previous argument that might be exploited to turn the gravitational scalar ϕ into the QCD axion. Inspired by the actual way that the field couples to fermions in EC gravity (see Appendix B of Supplemental Material [61] and Refs. [39,60,69,70,89,90] therein), let us generalize the interaction (7) as

$$\begin{aligned} \mathcal{L} \supset & \mathcal{Z}_V(\phi) \partial_\mu \phi J_V^\mu + \mathcal{Z}_A(\phi) \partial_\mu \phi J_A^\mu + \mathcal{Z}_{VV}(\phi) \frac{J_V^{\mu 2}}{M_P^2} \\ & + \mathcal{Z}_{AA}(\phi) \frac{J_A^{\mu 2}}{M_P^2} + \mathcal{Z}_{AV}(\phi) \frac{J_A^\mu J_V^\mu}{M_P^2}, \end{aligned} \quad (13)$$

where M_P is the Planck mass and $\mathcal{Z}(\phi)$ are *a priori* arbitrary functions of the dimensionless ϕ . The axial current J_A^μ was defined previously in (6), while

$$J_V^\mu = \bar{\Psi} \gamma^\mu \Psi \quad (14)$$

is its vectorial counterpart. The crucial point is that Eq. (13) does not feature any shift symmetry of ϕ , and so the restriction of our previous argument is evaded.

Perturbative effects—Since the factors $\mathcal{Z}$ in Eq. (13) depend on ϕ , these interactions induce anomalous Yukawa-type couplings between the $\mathcal{Z}(\phi)$ functions and quarks of the schematic form

$$\mathcal{L}_{\text{Yukawa}} \propto i f(\phi) \bar{\Psi}_L m_\Psi \Psi_R + \text{H.c.}, \quad (15)$$

where $f(\phi)$ controls the strength and precise form of the interaction; it is related to the various $\mathcal{Z}(\phi)$ ’s. Terms

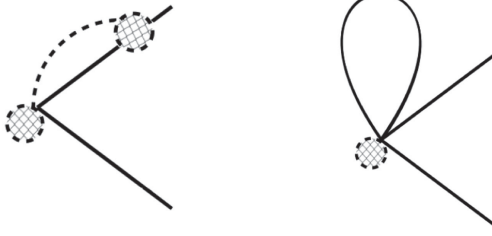


FIG. 1. Exemplary Feynman diagrams generating Yukawa-type couplings between the ϕ and quarks. We used a blob to denote ϕ insertions, which should be dressed appropriately depending on the $\mathcal{Z}$ functions that appear in the interactions. As usual, dashed and solid lines stand for scalar and fermionic fields, respectively. In order to obtain Yukawa couplings with γ_5 , the corresponding diagrams should involve both the axial J_A^μ and vectorial J_V^μ fermionic currents.

as in (15) appear already at the one-loop level, due to diagrams in Fig. 1. In this respect, what we encounter here is close in spirit to the Kim-Shifman-Vainshtein-Zakharov construction [19,20], where the axion couplings to the SM quarks are generated at the loop level. The left diagram in Fig. 1 gives a contribution to $f(\phi) \propto \mathcal{Z}_V(\phi)\mathcal{Z}_A(\phi)$, and the right graph $f(\phi) \propto \mathcal{Z}_{AV}(\phi)$. The coefficients in front of the $\mathcal{Z}$ factors are divergent and thus depend on the UV completion of the theory, which is power-counting nonrenormalizable.

We therefore see that ϕ can indeed communicate non-trivially with QCD, which is a necessary requirement for solving the strong CP problem by gravity. The induced effective potential for ϕ and $\bar{\theta}$ is given by Eq. (4) with $a/f_a \mapsto f(\phi)$, i.e.,

$$V = -m_\pi^2 f_\pi^2 \sqrt{1 - \frac{4m_u m_d}{(m_u + m_d)^2} \sin^2 \left(\frac{1}{2} [\bar{\theta} - f(\phi)] \right)}. \quad (16)$$

If V were the only contribution to the potential, its minimum at $\phi = f^{-1}(\bar{\theta})$ could lead to a dynamical vanishing of CP violation.

Several conditions must be satisfied for this solution to be viable. First, the domain of possible values of the function $f(\phi)$ should include the interval $[0, 2\pi)$ to compensate for any value of the θ angle, i.e., it should be $f(\phi) \sim 1$. Second, the QCD contribution to the axion potential should be dominant, i.e., the gravitationally induced axion mass must be much smaller than its QCD value. Besides a possible tree-level axion mass, the interactions in Eq. (13) generate the axion potential as well; see Fig. 2. They lead to two-loop contributions $\delta V \propto \mathcal{Z}_V(\phi)^2, \mathcal{Z}_A(\phi)^2, \mathcal{Z}_{VV}(\phi)$, and $\mathcal{Z}_{AA}(\phi)$. Similarly to the function $f(\phi)$, the coefficients in front of these factors depend on the UV completion of the theory. To solve the strong CP problem, we must have $\delta V \ll \Lambda_{\text{QCD}}^4$, where $\Lambda_{\text{QCD}} \sim 100$ MeV is the QCD scale.

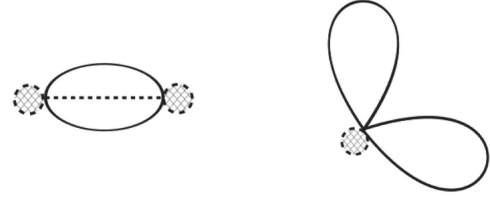


FIG. 2. Exemplary two-loop Feynman diagrams, originating from the couplings between the gravitational scalar and fermions in Eq. (13), that feed into the potential for ϕ .

Let us consider different possibilities for the UV completion. First, let us make a naive estimate of different divergent diagrams, assuming that they are all saturated at some universal UV cutoff Λ , much exceeding the electro-weak scale, which can be as large as the Planck scale. Then, omitting numerical coefficients $\mathcal{O}(1)$,

$$f(\phi) \sim \frac{\Lambda^2}{(4\pi f_a)^2} \left(\mathcal{Z}_V \mathcal{Z}_A + \mathcal{Z}_{AV} \frac{f_a^2}{M_P^2} \right), \quad (17)$$

and

$$\delta V \sim \frac{m_\psi^2 \Lambda^4}{(16\pi^2 f_a)^2} \left((\mathcal{Z}_A^2 + \mathcal{Z}_V^2) + (\mathcal{Z}_{AA} + \mathcal{Z}_{VV}) \frac{f_a^2}{M_P^2} \right). \quad (18)$$

Clearly, the requirement $f(\phi) \sim 1$ leads to $\delta V \gg \Lambda_{\text{QCD}}^4$. So, the conditions formulated in the previous paragraph cannot be met if the construction of the UV complete theory follows the naive power counting in the estimates of the Feynman graphs, unless different contributions to δV are severely fine-tuned.

Second, let us assume that the complete theory is such that all powerlike divergences are irrelevant [91,92], which can be implemented by using, e.g., dimensional regularization (DimReg). Then, an estimate of $f(\phi)$ and δV is given by the formulas [(17) and (18)] with the replacement $\Lambda \rightarrow m_\psi$. So, any communication with QCD due to the interactions (13) is practically nonexistent. In other words, if the only perturbative contributions were associated with the interactions (13), no solution to the strong CP problem would be achieved: roughly speaking, the nearly massless gravitational scalar field in this case is “orthogonal” to the required axion direction in the field space. In summary, we still have not managed to salvage the axion of gravitational origin.

Weyl-invariant R^2 gravity—Now we turn to the Weyl-invariant EC gravity as proposed in [39], and briefly presented across Appendixes B and C of Supplemental Material [93]. The self-consistency of the theory requires that quantum corrections be treated in a manner that preserves the Weyl symmetry. This singles out the use of the “scale-invariant” prescription [104,105] to regularize Feynman diagrams, where divergences are removed in a scale-invariant way. The simplest realization is based on

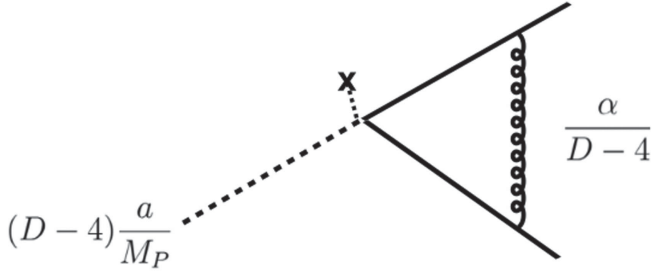


FIG. 3. Contribution to the axion-fermion coupling from evanescent operators. The factor $(D-4)$ multiplying a/M_P comes from the expansion of $(M_P^2 + M_P a)^{D-4}$ around $D=4$ and $a=0$. This is compensated by the divergent loop factor, leading to a finite contribution. As before, dashed and solid lines stand for scalar and fermionic fields, respectively; the wavy line denotes gluons, while $\times$ is the vacuum expectation value of the Higgs.

DimReg. The theory is extended to $D \neq 4$ dimensions, in a Weyl-invariant way, by multiplication of the terms in the classical Lagrangian with appropriate functions of the scalar fields present (Higgs h taken in the unitary gauge, dilaton χ that generates the Planck scale M_P , and axion a) [106]. For example, the Weyl-invariant Higgs-fermion interaction can be written as [107]

$$\int d^D x \sqrt{g} \bar{\Psi} \Psi h \chi^{D-4} F(a/\chi, h/\chi), \quad (19)$$

where $F \rightarrow 1$ when $D \rightarrow 4$, with a similar construction for the rest of the terms in the classical action. The finite radiative corrections are derived as in the 't Hooft minimal subtraction scheme by removing the poles in $D-4$. The results depend on the function F , which determines the UV completion of the theory. It cannot be fixed by a symmetry principle. What is important is that the effective action (in the Einstein frame) contains the higher-dimensional operators induced by these evanescent operators: the factors of $D-4$ in the interaction strength are compensated by the poles in $D-4$ appearing because of the loops [108]. As a result, the first and second terms that appear in Eq. (11) are generated, but now with uncorrelated numerical coefficients, depending on the function F . To illustrate this point, in Fig. 3 we have depicted how a Yukawa-type mixing can be generated at one loop. Consequently, the required coupling of the gravitational axion to these operators is already present in the Weyl-invariant EC theory, enabling the gravitational solution of the strong CP problem.

The discussion above was purely perturbative. The gravitational axion will remain a viable option only if non-perturbative gravity contributions to the axionic potential remain negligible.

Outlook—A particularly appealing candidate for the QCD axion is a scalar field arising purely from gravity. Such realizations are minimal in that the field is already

present in the gravitational sector and automatically couples to fermionic currents, and thus, naively, also to QCD via the chiral anomaly. This feature distinguishes them from more conventional axion models that rely on extra symmetries or fields.

We have shown, however, that previous arguments for how a QCD axion can be generated from gravity are ultimately flawed. Furthermore, addressing the shortcomings of these proposals appears to be significantly more challenging than one might have anticipated.

Concretely, our findings can be summarized as follows: (1) in theories where the would-be axion couples only derivatively to fermionic currents, its shift symmetry remains unbroken and so its presence is irrelevant for the strong CP problem. (2) If the coupling to fermions explicitly breaks shift symmetry, then quantum corrections can, in principle, open the communication with QCD. However, the *perturbative* interactions are either too weak to be viable or suffer from a severe fine-tuning problem. (3) These results leave one avenue for realizing a viable gravitational axion in the framework of the Weyl-invariant Einstein-Cartan gravity [39], where the coupling to QCD can arise from quantum Weyl invariance.

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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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