Indications against Dynamical *CPT* Symmetry Restoration in Quantum Gravity

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CPT symmetry is at the heart of the standard model of particle physics and experimentally very well tested, but expected to be broken in some approaches to quantum gravity. It thus becomes pertinent to explore which of the two alternatives is realized: (i) *CPT* symmetry is emergent, so that it is restored in the low-energy theory, even if it is broken beyond the Planck scale, (ii) *CPT* symmetry cannot be emergent and must be fundamental, so that any approach to quantum gravity, in which *CPT* is broken, is ruled out. We explore this by calculating the renormalization group flow of *CPT*-violating interactions under the impact of quantum fluctuations of the metric. We find that *CPT* symmetry cannot be emergent and conclude that quantum-gravity approaches must avoid the breaking of *CPT* symmetry. As a specific example, we discover that in asymptotically safe quantum gravity *CPT* symmetry remains intact, if it is imposed as a fundamental symmetry, but it is badly broken at low energies if a tiny amount of *CPT* violation is present in the transplanckian regime.

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CPT symmetry: Emergent or fundamental?—It is rare that experimental data can be used to constrain quantum-gravity effects, see Ref. [1] for notable exceptions. Here, we leverage strong constraints on *CPT* violation in the standard model (SM) to learn about the properties of quantum gravity at (trans) Planckian scales.

Symmetries are at the heart of particle physics. One of the most fundamental symmetries of particle physics is *CPT* symmetry, which combines the three discrete symmetries *C* (charge conjugation), *P* (parity transformation), and *T* (time reversal) into one symmetry operation.

CPT symmetry is experimentally very well tested [2–23]. This impressive experimental effort has to date not fully been appreciated regarding its implications for quantum gravity. Instead, on theoretical grounds it has been argued that quantum gravity effects could violate *CPT* symmetry, see, e.g., Refs. [1,24–37]. This may in fact be expected, if quantum spacetime is discrete, such that concepts of continuum quantum field theory (QFT) lose their meaning.

Against this background, an urgent question arises: if quantum gravity breaks *CPT*, why have the corresponding effects not been detected in experiments? There are two alternatives: (i) *CPT* symmetry is broken above the Planck scale, but restored at lower energies, i.e., it is an *emergent* symmetry, (ii) *CPT* symmetry cannot be emergent and must be imposed at the *fundamental* level. In this case any quantum-gravity theory that breaks *CPT* symmetry is ruled out.

A symmetry is emergent if all symmetry-violating terms are irrelevant under the Renormalization Group (RG) flow, i.e., they are driven to zero as one evolves the theory from high to low energy scales. A symmetry cannot be emergent, if symmetry-violating terms are RG relevant, because relevant couplings grow under the RG flow to low energies [38]. We can thus decide between the two alternatives by calculating loop corrections from quantum gravity to the leading-order *CPT*-violating terms.

CPT symmetry violation for standard model fields—*CPT* symmetry holds (under some assumptions) in any local, Lorentz invariant quantum field theory. Its violation in perturbatively renormalizable interactions is automatically tied to the violation of Lorentz invariance [39–42]. Loop corrections to various Lorentz-violating terms have been explored, see, e.g., Ref. [44] for matter loops and [45] for gravity loops and [46] for a summary.

We will analyze the quantum-gravity effects on the leading-order *CPT* violating terms for the matter fields in the standard model, i.e., for a scalar, fermion, and gauge field [47], which are [39,40]

$$S_{\text{scalar}}^{CPT-\text{odd}} = i g_\phi \int \mathrm{d}^4x \sqrt{g} n_\mu (\phi^* \partial^\mu \phi - \phi \partial^\mu \phi^*), \qquad (1)$$

$$S_{\text{photon}}^{CPT-\text{odd}} = \frac{1}{2} \int \mathrm{d}^4x \sqrt{g} g_\gamma n_\mu \epsilon^{\mu\nu\rho\sigma} A_\nu F_{\rho\sigma}, \qquad (2)$$

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$$S_{\mathrm{fermion}}^{CPT-\mathrm{odd}} = -\int \mathrm{d}^4x\sqrt{g}n_\mu(g_\psi\bar{\psi}\gamma^\mu\psi + h_\psi\bar{\psi}\gamma^\mu\gamma_5\psi). \tag{3}$$

Our notation separates a scalar coupling from a vector, such that it relates to the notation in Refs. [39,40] by $g_\phi n^\mu = (k_\phi)^\mu$, $g_\gamma n^\mu = (k_{AF})^\mu$, as well as $g_\psi n_\mu = a_\mu$, and $h_\psi n_\mu = b_\mu$. To explain their form, let us recall that the standard kinetic terms are *CPT* even and a derivative ∂_μ transforms into $-\partial_\mu$ under *PT* (and is invariant under *C*). Thus, to build a *CPT* odd term, we substitute one derivative in the kinetic term by an external vector n_μ which is *CPT* even. This external vector field violates Lorentz symmetry, because it singles out a preferred frame. The combinations $g_\phi n_\mu$, $g_\gamma n_\mu$, $g_\psi n_\mu$, and $h_\psi n_\mu$ parameterize the amount of *CPT* violation in the theory and are experimentally tightly constrained [48]. For photons, such a term results in birefringence [7], which is tightly constrained, e.g., from cosmic microwave background observations, to $g_\gamma n_\mu < 10^{-43}$ GeV [4,6,7], with more recent constraints [23] stronger by another order of magnitude. Similar constraints ($g_\gamma n_\mu < 10^{-41}$ GeV) arise from observation of photons that propagate over astrophysical distances [2]. For the Higgs field in the SM, constraints arise from lepton magnetic moments, and constrain $|g_\phi n_\mu|^2 < 10^{-29}$ GeV^2 [22]. In addition, already before the experimental discovery of the Higgs field, constraints were derived based on loop effects of the Higgs field, reading $g_\phi n_\mu < 10^{-29}$ GeV [3]. For fermions, constraints exist for the various quarks and leptons of the SM. For instance, for the τ, $|n_\mu h_\psi| < 10^{-10}$ GeV [17] and for the μ, $|n_\mu h_\psi| < 10^{-23}$ GeV [5]. For stable particles, Penning-trap experiments provide high-precision comparisons of charge-to-mass ratios of particles and antiparticles, resulting in limits on $n_\mu h_\psi < 10^{-17}$ GeV [19] and even $n_\mu h_\psi < 10^{-25}$ GeV [14,18]. For neutrinos, constraints from tritium decay amount to $n_\mu g_\psi < 10^{-5}$ GeV [21], and $n_\mu g_\psi < 10^{-20}$ GeV from neutrino oscillations [9–11,13,15,16]. In the quark sector, bounds of 10^{-18} GeV can be derived from Kaon oscillations [8,12,20]. Against this impressive experimental effort, the fate of *CPT* symmetry in quantum gravity becomes a pressing question.

Quantum-gravity fluctuations and CPT restoration—If quantum-gravity effects violate *CPT*, one would generically expect that the combinations $g_\phi n_\mu$ etc. are all $\mathcal{O}(M_{\mathrm{Planck}}) \approx 10^{19}$ GeV, which is clearly many orders of magnitude above the experimental bounds. This is a dimensional estimate which may oversimplify the actual behavior of the system. In particular, if couplings acquire anomalous scaling dimensions due to quantum fluctuations, their values at low energy may be very far from such dimensional estimates. Because the standard model is perturbative at energy scales below the Planck scale (with the exception of nonperturbative QCD effects which are irrelevant in the context of the above experimental bounds), quantum fluctuations of matter cannot induce anomalous scaling dimensions which are large enough to solve the tension between experimental bounds and theoretical expectation without fine-tuning. Indeed, this expectation is borne out in the explicit study in Ref. [49]. We therefore calculate the RG flow of the *CPT*-violating couplings under the impact of quantum fluctuations of gravity. These have, to the best of our knowledge, never been calculated before and are the last remaining way to avoid the requirement that the *CPT*-violating couplings have to be set to values extremely close to zero already in the UV.

There is a seeming ambiguity in this setting, which is the assignment of a scaling dimension to n_μ. For dynamical fields, scaling dimensions can be assigned by considering their kinetic terms in the action, which must be dimensionless. However, n_μ in our setting is not a dynamical field, but simply an externally given, fixed field. It does therefore not have a fixed scaling dimension and it is consistent to assign to it either scaling dimension 1, making the couplings $g_{\phi,\gamma,\psi}$ and h_ψ dimensionless; or assign to it dimension 0, making the couplings $g_{\phi,\gamma,\psi}$ and h_ψ dimensionful, with canonical scaling dimension 1 [50].

This ambiguity drops out when we consider the dependence of the combinations $g_i n_\mu$ ($i = \phi, \gamma, \psi$, $g_4 = h_\psi$) under the RG flow, because these always have dimension 1. For these combinations, it holds that

$$\beta_{g_i n_\mu} = k\partial_k(g_i n_\mu) = f_i g_i n_\mu + \ldots \tag{4}$$

The factors f_i parameterize the effect of gravitational fluctuations. We can solve this linear equation (assuming that $f_i \approx$ const). We set initial conditions for the RG flow at a scale Λ_{UV}, which is a UV cutoff scale beyond which our analysis ceases to be valid. For instance, for discrete quantum-gravity theories, Λ_{UV} would be associated with the discreteness scale. We assume $\Lambda_{\mathrm{UV}} \gtrsim M_{\mathrm{Planck}}$ and have

$$(g_i n_\mu)(k) = (g_i n_\mu)(\Lambda_{\mathrm{UV}}) \cdot \left(\frac{k}{\Lambda_{\mathrm{UV}}}\right)^{f_i}. \tag{5}$$

The canonical scaling dimension is absorbed into the initial condition and only the anomalous dimension f_i determines whether or not the coupling grows or increases towards the IR [52].

Thus, a discrepancy with the experimental result $g_i n_\mu \ll 10^{19}$ GeV can only be avoided (i) if the UV value of the coupling is extremely small, $(g_i n_\mu)_{\Lambda_{\mathrm{UV}}} \lesssim (g_i n_\mu)_{\mathrm{exp}}$ (assuming $f_i \approx 0$), or (ii) if the anomalous scaling dimensions f_i are large enough so that a large value of the coupling at the initial scale is multiplied by a suppression factor much smaller than one.

Possibility (i) may be realized, if a quantum-gravity theory has a mechanism that requires or at least allows for $g_i n_\mu(k \simeq M_{\mathrm{Pl}}) = 0$. We will comment on and explore two such possibilities, namely string theory and asymptotically safe gravity, below.

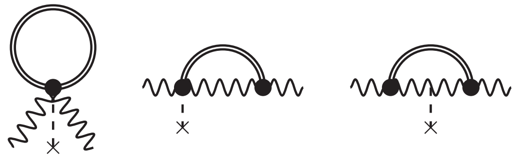

FIG. 1. Diagrams which contribute to the anomalous scaling dimension of the CPT-violating couplings. The double lines represent metric fluctuations, while the curly line represents any matter field, i.e., scalars, fermions, or gauge fields. The dashed line with cross indicates the insertion of the external vector field n_μ, which can sit either at the vertex or the propagator. The FRG regulator can be inserted on either of the internal lines of a given diagram.

Possibility (ii) may be realized if $f_i > 0$.

Metric fluctuations couple directly to the CPT-odd interaction terms in Eqs. (1)–(3) and generate diagrams of the form shown in Fig. 1, which we evaluate with functional RG (FRG) techniques, see Refs. [53–56], and Refs. [57–59] for reviews.

We approximate the gravitational dynamics by the Euclidean Einstein-Hilbert action (neglecting higher-order curvature terms)

$$S_{\mathrm{EH}} = \frac{1}{16\pi G_N}\int \mathrm{d}^4x\sqrt{g}(-R+2\Lambda), \tag{6}$$

where G_N and Λ are the Newton coupling and cosmological constant, respectively, which we will treat as input parameters. We parameterize metric fluctuations by the dimensionless counterpart of the Newton constant $G = G_N k^2$ and the dimensionless cosmological constant $\Lambda = \bar{\Lambda}k^{-2}$. Both of these quantities also depend on the RG scale k, such that their values that parametrize the metric fluctuations are to be understood as the UV values and not their IR values.

The key result of our FRG computation is the anomalous scaling dimensions for the CPT-violating terms under the impact of gravitational fluctuations,

$$f_\phi = 0 + \mathcal{O}(G^2), \tag{7}$$

$$f_\gamma = -G\frac{1-4\Lambda}{2\pi(1-2\Lambda)^2} \xrightarrow{\Lambda\to 0} -\frac{G}{2\pi}, \tag{8}$$

$$f_\psi = -G\frac{41-44\Lambda}{80\pi(1-2\Lambda)^2} \xrightarrow{\Lambda\to 0} -\frac{41}{80\pi}G, \tag{9}$$

where the result for the scalar field holds at leading order $[\mathcal{O}(G)]$, and we perform a next-to-leading-order study $[\mathcal{O}(G^2)]$ in Supplemental Material [61].

We thus find the following results, see also Fig. 2: (1) For gauge fields, $f_\gamma < 0$ in the gravitational parameter space, except for a tiny region. Thus, metric fluctuations generically *increase*, not decrease $n_\mu g_\gamma$, see Supplemental

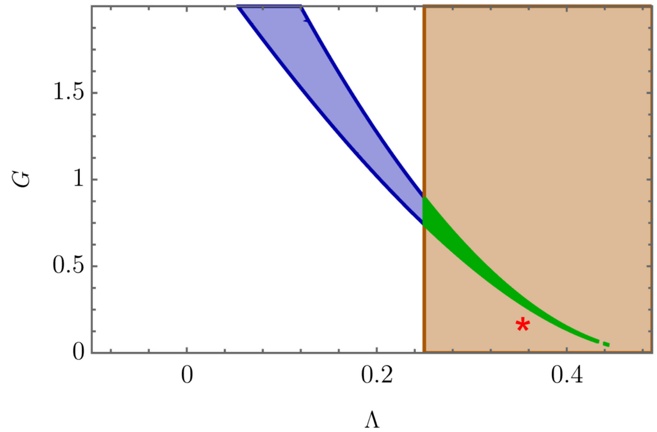

FIG. 2. Region in the gravitational parameter space, where $f_\gamma > 0$ (orange), and $f_\phi > 0$ (blue). We show the small overlap of both regions in green. In this region, CPT symmetry in the scalar and gauge sector is emergent. In contrast, violations of CPT symmetry in the fermion sector still grow towards low energies, even in the green region. The red star indicates the asymptotically safe fixed point for standard model matter content [60]. Its size should not be understood as an estimation of systematic uncertainties; those are expected to be larger.

Material for studies of the robustness [61]. (2) For fermions, $f_\psi < 0$ for both g_ψ and h_ψ everywhere in parameter space. Thus, metric fluctuations always *increase*, not decrease the fermionic CPT violating couplings. (3) For scalars, the leading-order contribution to f_ϕ vanishes. We therefore proceed to next-to-leading order, where we obtain a small contribution that is generically negative.

In summary, we do not find a region in parameter space where metric fluctuations drive the CPT-violating couplings towards zero. Instead, metric fluctuations even increase these couplings further, exacerbating the problem of reconciling experimental constraints with theoretical expectations. To demonstrate the severity of the required fine-tuning, we consider the observational constraint $g_\gamma n_\mu < 10^{-43}$ GeV [4,6,7] and confront it with our result for f_γ, setting $\Lambda = 0$ and $G = 1$ to obtain a concrete numerical example. For each order of magnitude in scales over which quantum-gravity fluctuations are dynamically important, and where we assume that we can approximate $f_\gamma = \mathrm{const}$, we obtain an enhancement factor of approximately 1.4. An enhancement is the opposite of the effect that would be needed; in fact, if CPT would be violated in the deep UV, but suppressed by gravitational fluctuations such that observational bounds were met, $f_i \gtrsim 10$ would be necessary to avoid fine-tuning. In contrast, we have $f_\gamma \approx -0.16$ for $G = 1$.

We conclude that, under the assumptions underlying our computation, any quantum theory of gravity that violates CPT and Lorentz symmetry can only be reconciled with observations at the cost of extreme fine-tuning of CPT-violating couplings to near-zero at Λ_{UV}. By adding higher-order CPT-violating operators, this conclusion becomes even stronger: to dynamically restore CPT symmetry at low energies, all of those operators need to be driven to zero by metric fluctuations. As soon as a single operator increases towards lower energies, CPT symmetry cannot be restored without fine-tuning.

The cost of prohibiting CPT violation in asymptotic safety—Asymptotically safe quantum gravity is a quantum field theory of the metric. Based on a proposal by Weinberg [72], it solves the perturbative nonrenormalizability of general relativity through an interacting fixed point of the RG. At such a fixed point, all interactions that correspond to perturbative counterterms of general relativity are generically present. They do not all correspond to independent free parameters, because only those interactions that form relevant perturbations of the fixed point introduce a free parameter. There is not just compelling evidence for the existence of such a fixed point, see Refs. [58,59,73] for reviews, but also for a total of roughly three free parameters [74–77]. In addition, long-standing open questions, like the existence of the fixed point in Lorentzian signature as well as the absence of unphysical degrees of freedom, ghosts, are being addressed now, see Refs. [78,79] and Refs. [80–85], respectively. There is also significant evidence for a fixed point when the SM and some of its extensions are coupled to gravity [86–91], with predictive power for the matter sector [92–96], giving rise to a potential UV completion for the standard model [97–99]. A key property of the asymptotically safe fixed point is its near-perturbative nature, see Refs. [74,76,100–105]. This implies that calculations that focus on the power-counting relevant and marginal interactions are already producing reliable results in many cases. In the present case, this suggests that there is a leading order of the gravitational contribution to f, which consists in gravitational loops; and higher, subleading orders, which are based on contributions of power-counting irrelevant couplings.

A second important property is that calculations consistently find the preservation of global symmetries, see, e.g., Refs. [106–113] (and references therein), contrary to the so-called no-global-symmetries conjecture [114–117]. This ensures that there is a fixed point at which the CPT-violating couplings that we consider vanish, because Eq. (4). This does not guarantee that they vanish at all scales. If they correspond to RG-relevant parameters, they can depart from the fixed-point value and acquire finite values in the IR. In such a case, it remains a consistent possibility that the couplings are zero, but it is not the only possibility.

In the context of asymptotic safety, the ambiguity regarding the dimensionality of the couplings matters. This is because free parameters are determined by RG relevant couplings. Here it matters whether the vector n^μ has vanishing or nonvanishing mass dimension. For the case of vanishing mass dimension, the anomalous scaling dimensions f_i must be larger than 1 in order to overwhelm the canonical scaling dimension of the coupling and turn the coupling into an irrelevant coupling that stays at the fixed point for all scales. In contrast, for the case of a nonvanishing mass dimension for n_μ, the canonical scaling dimension must only be larger than 0 for the coupling to be irrelevant and thus fixed to zero at all scales.

The most conservative assumption is that n_μ has vanishing mass dimension and thus the f's must overwhelm a canonical scaling of -1 in order for the CPT-violating couplings to be irrelevant. Making CPT violation irrelevant in the U(1) gauge sector comes at the cost of a Landau-pole problem in the U(1) sector, at least in the current setup. This is because asymptotically safe gravity can solve the Landau-pole problem, if quantum gravity fluctuations antiscreen the $U(1)$ gauge coupling. There is overwhelming evidence that this happens [118] in most of the gravitational parameter space [93,95,120–124]. The very small remaining region in which the Landau-pole problem is not resolved is precisely the region in which the CPT-violating coupling becomes irrelevant [125], see Fig. 2, where we also indicate the location of the asymptotically safe fixed point. However, Christiansen *et al.* [126] have shown that a calculation that additionally accounts for the momentum dependence of the gravitational correction removes this remaining region, thus resolving the Landau-pole problem everywhere and rendering CPT-violation relevant everywhere.

We thus conclude that asymptotically safe gravity likely contains both CPT symmetric theories and CPT-violating theories in its landscape. CPT symmetry remains intact, if the relevant, CPT-violating perturbations of the fixed point are not used in constructing RG trajectories. In other words, it is consistent to *impose* CPT symmetry in asymptotically safe gravity, because it is not generated if CPT violation is set to zero exactly in the UV. However, we do not find a dynamical mechanism that would *explain* why the low-energy theory is CPT symmetric, because this would require CPT-violating interactions to be irrelevant at the fixed point.

Conclusions and outlook—We have discovered that breaking CPT symmetry at the level of dimension-three operators in quantum gravity is only compatible with experimental bounds at the cost of extreme fine-tuning. The already existing problem that a dimension-three operator comes with a coupling of the order of the Planck mass is exacerbated further by quantum-gravity fluctuations *reducing* the scaling dimension of the operator further and thus driving up the scaling dimension of the coupling. Our calculations are subject to approximations, e.g., to leading-order treatment of the gravitational sector. We check the robustness of these approximations in Supplemental Material [61]. In addition, we work in Euclidean signature, but expect that our results carry over to Lorentzian signature: within the effective-field-theory setup, gravitational physics is perturbative and the status of an analytical continuation is exactly the same as for nongravitational QFT; within the asymptotic-safety setup explicit Lorentzian calculations support close agreement with Euclidean ones [78,79,83,85,91,127–129], see also Refs. [43,88,130–132] for steps towards Lorentzian signature.

Therefore, we conclude that quantum-gravity theories must find a mechanism to preserve *CPT* symmetry exactly. Examples of such theories exist. For instance, in string theory, *CPT* symmetry is gauged [133] and must therefore—barring anomalies rendering the theory inconsistent [134]—stay an exact symmetry, even though it may be broken spontaneously [26]. More broadly, the no-global-symmetries conjecture in quantum gravity has been used to argue for the gauging of *CPT* [135]. We have discovered a second example here, by finding that in asymptotically safe quantum gravity, a fixed point of the *CPT*-symmetry-violating terms must lie at vanishing couplings. It is thus consistent to set these couplings exactly to zero at all scales.

This is not the only possibility consistent with a fixed point, because the *CPT*-violating terms are RG relevant and can thus depart from the fixed point. There is therefore no predictive power for *CPT* symmetry in asymptotic safety. The landscape of asymptotically safe theories may include theories with *CPT* violation.

Our results imply—within the technical limitations of our Letter—that any quantum gravity theory which satisfies our assumptions and violates *CPT* symmetry is likely not phenomenologically viable. This may in particular include theories in which spacetime is discrete and in which therefore P and T are likely not realized.

Finally, our results are also relevant in the context of proposals for cosmology which do not specify the details of the quantum-gravity era, but are based on *CPT* symmetry [136].

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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.

[1] A. Addazi *et al.*, Prog. Part. Nucl. Phys. **125**, 103948 (2022).
[2] B. Nodland and J. P. Ralston, Phys. Rev. Lett. **78**, 3043 (1997).
[3] D. L. Anderson, M. Sher, and I. Turan, Phys. Rev. D **70**, 016001 (2004).
[4] V. A. Kostelecky and M. Mewes, Phys. Rev. Lett. **99**, 011601 (2007).
[5] G. W. Bennett *et al.* (Muon (g-2) Collaboration), Phys. Rev. Lett. **100**, 091602 (2008).
[6] V. A. Kostelecky and M. Mewes, Astrophys. J. Lett. **689**, L1 (2008).
[7] V. A. Kostelecky and M. Mewes, Phys. Rev. D **80**, 015020 (2009).
[8] A. Di Domenico *et al.* (KLOE Collaboration), Found. Phys. **40**, 852 (2010).
[9] A. A. Aguilar-Arevalo *et al.* (MiniBooNE Collaboration), Phys. Lett. B **718**, 1303 (2013).
[10] Y. Abe *et al.* (Double Chooz Collaboration), Phys. Rev. D **86**, 112009 (2012).
[11] T. Katori (MiniBooNE Collaboration), Mod. Phys. Lett. A **27**, 1230024 (2012).
[12] D. Babusci *et al.* (KLOE-2 Collaboration), Phys. Lett. B **730**, 89 (2014).
[13] K. Abe *et al.* (Super-Kamiokande Collaboration), Phys. Rev. D **91**, 052003 (2015).
[14] Y. Ding and V. A. Kostelecký, Phys. Rev. D **94**, 056008 (2016).
[15] K. Abe *et al.* (T2K Collaboration), Phys. Rev. D **95**, 111101 (2017).
[16] B. Aharmim *et al.* (SNO Collaboration), Phys. Rev. D **98**, 112013 (2018).
[17] C. A. Escobar, J. P. Noordmans, and R. Potting, Phys. Rev. D **97**, 115030 (2018).
[18] Y. Ding, Symmetry **11**, 1220 (2019).
[19] Y. Ding and M. F. Rawnak, Phys. Rev. D **102**, 056009 (2020).
[20] A. Di Domenico, Symmetry **12**, 2063 (2020).
[21] M. Aker *et al.* (KATRIN Collaboration), Phys. Rev. D **107**, 082005 (2023).
[22] P. L. Cástulo, J. J. Toscano, and E. S. Tututi, Int. J. Mod. Phys. A **37**, 2250138 (2022).
[23] L. Caloni, S. Giardiello, M. Lembo, M. Gerbino, G. Gubitosi, M. Lattanzi, and L. Pagano, J. Cosmol. Astropart. Phys. 03 (2023) 018.
[24] R. M. Wald, Phys. Rev. D **21**, 2742 (1980).
[25] S. W. Hawking, Phys. Rev. D **32**, 2489 (1985).
[26] V. A. Kostelecky and R. Potting, Nucl. Phys. **B359**, 545 (1991).
[27] A. Strominger, Phys. Rev. D **48**, 5769 (1993).
[28] V. A. Kostelecky and R. Potting, Phys. Lett. B **381**, 89 (1996).
[29] J. R. Ellis, N. E. Mavromatos, and D. V. Nanopoulos, in *Workshop on K Physics* (Ed. Frontières, Gif-sur-Yvette, 1996), https://inspirehep.net/literature/421198.
[30] N. E. Mavromatos, Lect. Notes Phys. **669**, 245 (2005).
[31] N. E. Mavromatos, Found. Phys. **40**, 917 (2010).
[32] N. E. Mavromatos and S. Sarkar, Eur. Phys. J. C **73**, 2359 (2013).
[33] J. Ellis, N. E. Mavromatos, and S. Sarkar, Phys. Lett. B **725**, 407 (2013).

[34] M. Arzano and J. Kowalski-Glikman, Phys. Lett. B **760**, 69 (2016).
[35] B. Valeixo Bento, Causal Set Dynamics and The Problem of Time, Ph.D. thesis, Imperial Collage, London, 2018.
[36] M. Arzano, A. Bevilacqua, J. Kowalski-Glikman, G. Rosati, and J. Unger, Phys. Rev. D **103**, 106015 (2021).
[37] N. E. Mavromatos, Lect. Notes Phys. **1017**, 3 (2023).
[38] One can still avoid experimental bounds in this scenario by accepting an extreme fine-tuning to tiny values in the transplanckian regime.
[39] D. Colladay and V. A. Kostelecky, Phys. Rev. D **55**, 6760 (1997).
[40] D. Colladay and V. A. Kostelecky, Phys. Rev. D **58**, 116002 (1998).
[41] O. W. Greenberg, Phys. Rev. Lett. **89**, 231602 (2002).
[42] In quantum gravity, at best local Lorentz invariance holds. However, in a regime in which the theory is well described by small metric fluctuations about a Minkowski background, global Lorentz invariance is in principle recoverable. However, Lorentz symmetry is not dynamically restored in simple settings [43].
[43] B. Knorr, Phys. Lett. B **792**, 142 (2019).
[44] A. I. Hernández-Juárez, J. Montaño, H. Novales-Sánchez, M. Salinas, J. J. Toscano, and O. Vázquez-Hernández, Phys. Rev. D **99**, 013002 (2019).
[45] A. Eichhorn, A. Platania, and M. Schiffer, Phys. Rev. D **102**, 026007 (2020).
[46] A. F. Ferrari, J. R. Nascimento, and A. Y. Petrov, Eur. Phys. J. C **80**, 459 (2020).
[47] We work in an extension of the standard model, where right-handed neutrinos are added, such that all fermion fields of the standard model can be combined into Dirac spinors, once we ignore the interactions of the standard model, as is sufficient for our analysis.
[48] Often, constraints are reported on individual components of these interactions; because we are only interested in the order of magnitude of the strongest constraints, we do not make this distinction here. Hence, if we quote bounds on vectors, we refer to the strongest bound on any component of that vector; all details can be found in the corresponding papers.
[49] D. Colladay and P. McDonald, Phys. Rev. D **77**, 085006 (2008).
[50] In settings in which the vector is dynamical, e.g., in Einstein-aether-theory [51], a dimension can be assigned.
[51] T. Jacobson, Proc. Sci. QG-PH2007 (**2007**) 020 [arXiv: 0801.1547].
[52] We will see below that the situation is slightly more subtle in the case of an asymptotically safe RG fixed point.
[53] C. Wetterich, Phys. Lett. B **301**, 90 (1993).
[54] T. R. Morris, Int. J. Mod. Phys. A **09**, 2411 (1994).
[55] U. Ellwanger, Z. Phys. C **62**, 503 (1994).
[56] M. Reuter, Phys. Rev. D **57**, 971 (1998).
[57] N. Dupuis, L. Canet, A. Eichhorn, W. Metzner, J. M. Pawlowski, M. Tissier, and N. Wschebor, Phys. Rep. **910**, 1 (2021).
[58] F. Saueressig, *The Functional Renormalization Group in Quantum Gravity* (Springer, Singapore, 2023).
[59] J. M. Pawlowski and M. Reichert, *Quantum Gravity from Dynamical Metric Fluctuations* (Springer, Singapore, 2024).
[60] N. Christiansen, K. Falls, J. M. Pawlowski, and M. Reichert, Phys. Rev. D **97**, 046007 (2018).
[61] See Supplemental Material at http://link.aps.org/supplemental/10.1103/frd8-pg3d for additional details on the methodology and setup, and for various tests for robustness of our results; includes Refs. [62–71].
[62] J. Berges, N. Tetradis, and C. Wetterich, Phys. Rep. **363**, 223 (2002).
[63] J. M. Pawlowski, Ann. Phys. (Amsterdam) **322**, 2831 (2007).
[64] H. Gies, Lect. Notes Phys. **852**, 287 (2012).
[65] B. Delamotte, Lect. Notes Phys. **852**, 49 (2012).
[66] O. J. Rosten, Phys. Rep. **511**, 177 (2012).
[67] J. Braun, J. Phys. G **39**, 033001 (2012).
[68] M. Reuter and F. Saueressig, New J. Phys. **14**, 055022 (2012).
[69] M. Reichert, Proc. Sci. 384 (**2020**) 005.
[70] A. Eichhorn, Phys. Rev. D **86**, 105021 (2012).
[71] G. P. de Brito, B. Knorr, and M. Schiffer, Phys. Rev. D **108**, 026004 (2023).
[72] S. Weinberg, Ultraviolet divergences in quantum theories of gravitation, in *General Relativity: An Einstein Centenary Survey* edited by S. W. Hawking and W. Israel (Cambridge University Press, 1980), pp. 790–831, https://www.cambridge.org/ag/universitypress/subjects/physics/cosmology-relativity-and-gravitation/general-relativity-einstein-centenary-survey?format=PB.
[73] A. Eichhorn and M. Schiffer, arXiv:2212.07456.
[74] K. Falls, D. F. Litim, K. Nikolakopoulos, and C. Rahmede, Phys. Rev. D **93**, 104022 (2016).
[75] T. Denz, J. M. Pawlowski, and M. Reichert, Eur. Phys. J. C **78**, 336 (2018).
[76] K. Falls, C. R. King, D. F. Litim, K. Nikolakopoulos, and C. Rahmede, Phys. Rev. D **97**, 086006 (2018).
[77] K. Falls, N. Ohta, and R. Percacci, Phys. Lett. B **810**, 135773 (2020).
[78] E. D'Angelo, Phys. Rev. D **109**, 066012 (2024).
[79] F. Saueressig and J. Wang, Phys. Rev. D **111**, 106007 (2025).
[80] B. Knorr, C. Ripken, and F. Saueressig, Classical Quantum Gravity **36**, 234001 (2019).
[81] A. Platania and C. Wetterich, Phys. Lett. B **811**, 135911 (2020).
[82] B. Knorr and M. Schiffer, Universe **7**, 216 (2021).
[83] J. Fehre, D. F. Litim, J. M. Pawlowski, and M. Reichert, Phys. Rev. Lett. **130**, 081501 (2023).
[84] A. Platania, J. High Energy Phys. 09 (2022) 167.
[85] A. Pastor-Gutiérrez, J. M. Pawlowski, M. Reichert, and G. Ruisi, Phys. Rev. D **111**, 106005 (2025).
[86] P. Donà, A. Eichhorn, and R. Percacci, Phys. Rev. D **89**, 084035 (2014).
[87] J. Meibohm, J. M. Pawlowski, and M. Reichert, Phys. Rev. D **93**, 084035 (2016).
[88] J. Biemans, A. Platania, and F. Saueressig, J. High Energy Phys. 05 (2017) 093.
[89] N. Alkofer and F. Saueressig, Ann. Phys. (Amsterdam) **396**, 173 (2018).
[90] C. Wetterich and M. Yamada, Phys. Rev. D **100**, 066017 (2019).

[91] G. Korver, F. Saueressig, and J. Wang, Phys. Lett. B **855**, 138789 (2024).
[92] M. Shaposhnikov and C. Wetterich, Phys. Lett. B **683**, 196 (2010).
[93] U. Harst and M. Reuter, J. High Energy Phys. 05 (2011) 119.
[94] A. Eichhorn and A. Held, Phys. Lett. B **777**, 217 (2018).
[95] A. Eichhorn and F. Versteegen, J. High Energy Phys. 01 (2018) 030.
[96] A. Eichhorn and A. Held, Phys. Rev. Lett. **121**, 151302 (2018).
[97] R. Alkofer, A. Eichhorn, A. Held, C. M. Nieto, R. Percacci, and M. Schröfl, Ann. Phys. (Amsterdam) **421**, 168282 (2020).
[98] K. Kowalska, S. Pramanick, and E. M. Sessolo, J. High Energy Phys. 08 (2022) 262.
[99] A. Pastor-Gutiérrez, J. M. Pawlowski, and M. Reichert, SciPost Phys. **15**, 105 (2023).
[100] K. Falls, D. F. Litim, K. Nikolakopoulos, and C. Rahmede, arXiv:1301.4191.
[101] K. G. Falls, D. F. Litim, and J. Schröder, Phys. Rev. D **99**, 126015 (2019).
[102] A. Eichhorn, P. Labus, J. M. Pawlowski, and M. Reichert, SciPost Phys. **5**, 031 (2018).
[103] A. Eichhorn, S. Lippoldt, J. M. Pawlowski, M. Reichert, and M. Schiffer, Phys. Lett. B **792**, 310 (2019).
[104] A. Eichhorn, S. Lippoldt, and M. Schiffer, Phys. Rev. D **99**, 086002 (2019).
[105] Y. Kluth and D. F. Litim, Phys. Rev. D **108**, 026005 (2023).
[106] G. Narain and R. Percacci, Classical Quantum Gravity **27**, 075001 (2010).
[107] A. Eichhorn and H. Gies, New J. Phys. **13**, 125012 (2011).
[108] P. Labus, R. Percacci, and G. P. Vacca, Phys. Lett. B **753**, 274 (2016).
[109] A. Eichhorn and S. Lippoldt, Phys. Lett. B **767**, 142 (2017).
[110] J. Meibohm and J. M. Pawlowski, Eur. Phys. J. C **76**, 285 (2016).
[111] G. P. de Brito, A. Eichhorn, and R. R. L. d. Santos, J. High Energy Phys. 11 (2021) 110.
[112] C. Laporte, A. D. Pereira, F. Saueressig, and J. Wang, J. High Energy Phys. 12 (2021) 001.
[113] A. Eichhorn, J. H. Kwapisz, and M. Schiffer, Phys. Rev. D **105**, 106022 (2022).
[114] T. Banks and L. J. Dixon, Nucl. Phys. **B307**, 93 (1988).
[115] T. Banks and N. Seiberg, Phys. Rev. D **83**, 084019 (2011).
[116] D. Harlow and H. Ooguri, Commun. Math. Phys. **383**, 1669 (2021).
[117] T. Daus, A. Hebecker, S. Leonhardt, and J. March-Russell, Nucl. Phys. **B960**, 115167 (2020).
[118] A seeming discrepancy with perturbative calculations has been addressed in [119].
[119] G. P. de Brito and A. Eichhorn, Eur. Phys. J. C **83**, 161 (2023).
[120] J.-E. Daum, U. Harst, and M. Reuter, J. High Energy Phys. 01 (2010) 084.
[121] S. Folkerts, D. F. Litim, and J. M. Pawlowski, Phys. Lett. B **709**, 234 (2012).
[122] N. Christiansen and A. Eichhorn, Phys. Lett. B **770**, 154 (2017).
[123] A. Eichhorn and M. Schiffer, Phys. Lett. B **793**, 383 (2019).
[124] G. P. De Brito, A. Eichhorn, and A. D. Pereira, J. High Energy Phys. 09 (2019) 100.
[125] The complementarity of the regions is not an accident, but has a structural reason: the beta function of the *CPT*-violating coupling receives no contributions that renormalize the vertex, because these cancel exactly among each other. The only remaining contribution is to the wave-function renormalization of the gauge field, which, by gauge symmetry, determines the renormalization of the gauge coupling.
[126] N. Christiansen, D. F. Litim, J. M. Pawlowski, and M. Reichert, Phys. Rev. D **97**, 106012 (2018).
[127] E. Manrique, S. Rechenberger, and F. Saueressig, Phys. Rev. Lett. **106**, 251302 (2011).
[128] E. D'Angelo, R. Ferrero, and M. B. Fröb, Classical Quantum Gravity **42**, 125008 (2025).
[129] J. M. Pawlowski, M. Reichert, and J. Wessely, arXiv:2507.22169.
[130] J. Biemans, A. Platania, and F. Saueressig, Phys. Rev. D **95**, 086013 (2017).
[131] A. Bonanno, T. Denz, J. M. Pawlowski, and M. Reichert, SciPost Phys. **12**, 001 (2022).
[132] F. Saueressig and J. Wang, J. High Energy Phys. 09 (2023) 064.
[133] J. Polchinski, *String Theory. Vol. 2: Superstring Theory and Beyond*, Cambridge Monographs on Mathematical Physics (Cambridge University Press, Cambridge, England, 2007).
[134] M. Dine and M. Graesser, J. High Energy Phys. 01 (2005) 038.
[135] D. Harlow and T. Numasawa, J. High Energy Phys. 01 (2026) 098.
[136] L. Boyle, K. Finn, and N. Turok, Phys. Rev. Lett. **121**, 251301 (2018).