Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

Poincaré Breaking and Gauge Invariance: A Road to Emergent Gravity and New Particles

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Abstract. In this talk, we discuss how gauge symmetries broken explicitly by a Poincare-breaking UV cutoff can be restored. We show that gauge symmetries can be restored by the introduction of affine curvature in reminiscence to the Higgs field. In fact, gauge symmetries get restored and general relativity emerges at the extremum of the metric-affine action. As per this point, we show emergence of the general relativity, reveal how its parameters relate to the flat spacetime loops, elucidate the new particle spectrum it brings along, and discuss its salient signatures. We show that the resulting field-theoretic plus gravitational setup can be probed via various phenomena ranging from collider experiments to black holes.

1. Introduction

There are scales in nature that are not particle masses. The confinement scale in QCD and the fundamental scale of gravity are such scales. In general, if the UV cutoff $\Lambda_\wp$ is not the mass of an elementary particle then it breaks Poincare symmetry as it is not a Casimir invariant of the Poincare group [1]. The simplest UV cutoff as such is the Lorentz-invariant but translation-breaking cut $-\Lambda_\wp^2 \leq \ell_\mu \ell^\mu \leq \Lambda_\wp^2$ on the loop momenta ℓ_μ. Under this cut, the effective QFT remains Lorentz-invariant but breaks explicitly all the gauge symmetries since each gauge boson acquires a mass proportional to $\Lambda_\wp$ [2, 3, 4, 5]. In addition, scalar mass-squareds (vacuum energy) develop quadratic (quadratic and quartic) sensitivity to the UV cutoff $\Lambda_\wp$ [6, 7].

In this talk, we will discuss how gauge symmetries broken explicitly by a Poincare-breaking UV cutoff can be restored by a mechanism reminiscent of the usual Higgs mechanism. Indeed, gauge symmetries associated with a massive vector boson are restored (or realized) by the introduction of the Higgs field [8, 9]. In reminiscent of this, we find that gauge symmetries broken explicitly by a Poincare-breaking UV cutoff can be restored by the introduction of the affine curvature [10, 11, 12]. In the former, what is central is the Higgs potential. In the latter, however, the central object iss the metric-affine action [13, 14, 15, 16].

2. Effective QFT

For the purpose of analyzing the power-law quantum divergences in isolation, it is desirable to detach them from the logarithmic divergences in terms of the scales they involve. This scheme of regularization, the so-called detached regularization [17], extends the usual dimensional regularization [18, 19, 20] to QFTs with a UV cutoff. In detached regularization, the effective action capturing physics of the full quantum action at low energies takes the form [17]

$$S_{eff}[\eta, F; \Lambda_\wp^2, \log\mu] = S_{tree}[\eta, F] + \delta S_{log}[\eta, \log\mu, F] + \delta S_{pow}[\eta, F; \Lambda_\wp^2, \log\mu] \quad (1)$$

Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

in which $S_{tree}[\eta, \digamma]$ stands for the tree-level QFT action and $\delta S_{log}[\eta, \log\mu, \digamma]$ collects the logarithmic loop corrections such that both actions are gauge-invariant and independent of the UV cutoff $\Lambda_\wp$. The last piece

$$\delta S_{pow}[\eta, \digamma; \Lambda_\wp^2, \log\mu] \;=\; \int d^4x\sqrt{-\eta}\left\{-c_O\Lambda_\wp^4 - \mathscr{M}^2\Lambda_\wp^2 - c_\phi\Lambda_\wp^2\phi^\dagger\phi + c_V\Lambda_\wp^2 \mathrm{tr}\left[V_\mu V^\mu\right]\right\} \quad (2)$$

comprises the power-law corrections. The loop factors c_ϕ and c_V depend on the QFT under concern but c_O and $\mathscr{M}^2$ are rather general at one loop [17]

$$c_O \;=\; c_O(\log\mu) \xrightarrow{\text{one loop}} \frac{(n_b - n_f)}{64\pi^2}, \quad (3)$$

$$\mathscr{M}^2 \;=\; \mathscr{M}^2(\log\mu) \xrightarrow{\text{one loop}} -\frac{1}{64\pi^2}\mathrm{str}\left[M^2 \log\frac{M^2}{\mu^2}\right], \quad (4)$$

in which $n_b(n_f)$ is the total number of bosons (fermions) and M^2 is the mass-squared matrix of the QFT fields. (Details of the detached regularization and its applications can be found in [17].)

3. Effective QFT in Curved Spacetime

QFTs are specific to the flat spacetime due to the necessity of Poincare symmetry for defining the notion of particle [21, 22]. They cannot be carried into curved spacetime because curved spacetime does not allow special states like the vacuum and detectable structures like the particles [23, 24].

Effective QFTs are different. They are the QFTs of long-wavelength quantum fields. They are essent,ally the classical theory improved with quantum corrections order by order in perturbation theory. In this regard, they are like the classical field theories and have therefore a natural affinity with the classical curved spacetime. In other words, it should be natural to carry effective QFTs into curved spacetime [11, 12]. In this regard, it is necessary to carry first the flat spacetime effective QFT (with metric $\eta_{\mu\nu}$ and the partial derivative ∂_μ) to curved spacetime (with metric $g_{\mu\nu}$ and covariant derivative ∇_μ of the Levi-Civita connection ${}^g\Gamma^\lambda_{\mu\nu}$). But, for the metric $g_{\mu\nu}$ to be curved, the effective QFT in curved spacetime must involve curvature of $g_{\mu\nu}$ (like, for example, the Ricci curvature $R_{\mu\nu}({}^g\Gamma)$). It can be tempting to add the requisite curvature terms by hand (Einstein-Hilbert term $M_0^2 g^{\mu\nu}R_{\mu\nu}({}^g\Gamma)$ plus higher-curvature terms) but extension of the effective QFT by such bare terms is fundamentally inconsistent because they come to mean that the curvature sector was left unrenormalized while the QFT sector was renormalized. Similar arguments hold also for curvature terms constructed with the bare QFT parameters in $S_{tree}[\eta, \digamma]$. In the face of these inconsistencies, one concludes that curvature must arise only in the loop-induced terms in the effective QFT [12]. To this end, definition of the Ricci curvature $[\nabla_\lambda, \nabla_\nu]V^\lambda = R_{\mu\nu}({}^g\Gamma)V^\mu$ reveals that the effective QFT in curved spacetime can develop curvature only in the gauge sector. But the gauge kinetic term $V_{\mu\nu}V^{\mu\nu}$ (with the field strength tensor $V_{\mu\nu} = \mathcal{D}_\mu V_\nu - \mathcal{D}_\nu V_\mu$ and gauge covariant derivative $\mathcal{D}_\mu$) can generate two-derivative structures like $[\nabla_\lambda, \nabla_\nu]V^\lambda$ via only by-parts integration. Indeed, $V_{\mu\nu}V^{\mu\nu}$ can be brought into the two-derivative form $V_\mu(\mathcal{D}^2)^{\mu\nu}V_\nu$ by applying by-parts integration. But by-parts integration cannot be a real source of curvature since after all $V_\mu(\mathcal{D}^2)^{\mu\nu}V_\nu$ is equivalent to $V_{\mu\nu}V^{\mu\nu}$. This observation implies that curvature must arise while passing from flat spacetime to the curved spacetime. In this regard, one notices that the kinetic construct [12]

$$I_V[\eta] \;=\; \int d^4x\sqrt{-\eta}\frac{1}{2}\mathrm{tr}[V_{\mu\nu}V^{\mu\nu}] \quad (5)$$

Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

is equivalent to

$$\bar{I}_V[\eta] \quad = \quad \int d^4x\sqrt{-\eta}\,\mathrm{tr}\big[V^\mu(-D^2\eta_{\mu\nu} + D_\mu D_\nu + iV_{\mu\nu})V^\nu + \partial_\mu(V_\nu V^{\mu\nu})\big] \tag{6}$$

under by-parts integration, with the flat spacetime gauge covariant derivative D_μ. Their difference vanishes identically

$$\Delta I_V[\eta] = -I_V[\eta] + \bar{I}_V[\eta] = 0 \tag{7}$$

with the general boundary term $\partial_\mu(V_\nu V^{\mu\nu})$ in $\bar{I}_V[\eta]$. But this very vanishing difference becomes non-vanishing

$$\Delta I_V[g] = -I_V[g] + \bar{I}_V[g] = -\int d^4x\sqrt{-g}\,\mathrm{tr}[V^\mu R_{\mu\nu}({}^g\Gamma)V^\nu] \tag{8}$$

when $I_V[\eta]$ and $\bar{I}_V[\eta]$ are taken to curved spacetime via the general covariance ($\eta_{\mu\nu} \to g_{\mu\nu}$ and $\partial_\mu \to \nabla_\mu$ so that $D_\mu \to \mathcal{D}_\mu$). The reason for this non-vanishing $\Delta I_V[g]$ is the aforementioned commutator $[\mathcal{D}_\lambda, \mathcal{D}_\nu]V^\lambda = (R_{\mu\nu}({}^g\Gamma) + iV_{\mu\nu})V^\mu$ in curved spacetime.

It proves useful to combine ΔI_V with the loop-induced gauge boson mass term

$$\delta S_V[\eta, \Lambda_\wp^2] = \int d^4x\sqrt{-\eta}\, c_V\Lambda_\wp^2\eta_{\mu\nu}\mathrm{tr}\left[V^\mu V^\nu\right] \tag{9}$$

in the power-law corrections in (2). It is clear that $\delta S_V[\eta, \Lambda_\wp^2]$ can be rewritten as

$$\delta\overline{S}_V[\eta, \Lambda_\wp^2] = \delta S_V[\eta, \Lambda_\wp^2] + c_V\Delta I_V[\eta] = S_V[\eta, \Lambda_\wp^2] \tag{10}$$

thanks to the fact that you $\Delta I_V[\eta] = 0$ in flat spacetime. This identity undergoes a nontrivial change [11, 12]

$$\delta\overline{S}_V[g, \Lambda_\wp^2, R] = \delta S_V[g, \Lambda_\wp^2] + c_V\Delta I_V[g] = \int d^4x\sqrt{-g}\, c_V\mathrm{tr}\left[V^\mu\left(\Lambda_\wp^2 g_{\mu\nu} - R_{\mu\nu}({}^g\Gamma)\right)V^\nu\right] \tag{11}$$

in curved spacetime in which $\Delta I_V[g] \neq 0$ as given in the relation (8). It is worth emphasizing that $\delta\overline{S}_V[g, \Lambda_\wp^2, R]$ is the only piece that explicitly involves curvature in the entire effective QFT in curved spacetime. In other words, without $\delta\overline{S}_V[g, \Lambda_\wp^2, R]$ (more precisely $\Delta I_V[g]$), the metric $g_{\mu\nu}$ would remain flat as it would have no kinetic term at all. As a result, the effective QFT action takes the form

$$S_{eff}[g, F; \Lambda_\wp^2, \log\mu, R] = S_{tree}[g, F] + \delta S_{log}[g, \log\mu, F] + \delta\overline{S}_{pow}[\eta, F; \Lambda_\wp^2, \log\mu, R] \tag{12}$$

in the curved spacetime of the metric $g_{\mu\nu}$. Its power-law part

$$\begin{aligned}\delta\overline{S}_{pow}[g, F; \Lambda_\wp^2, \log\mu, R] = \int d^4x\sqrt{-g}\Big\{ & - \quad c_O\Lambda_\wp^4 - \mathscr{M}^2\Lambda_\wp^2 - c_\phi\Lambda_\wp^2\phi^\dagger\phi \\ & + \quad c_V\mathrm{tr}\left[V^\mu\left(\Lambda_\wp^2 g_{\mu\nu} - R_{\mu\nu}({}^g\Gamma)\right)V^\nu\right]\Big\}\end{aligned} \tag{13}$$

is obtained by taking the power-law piece in (2) to curved spacetime via the relation (8).

It is clear that the entire effective QFT in curved spacetime (12) develops a curvature-dependent term only in the gauge sector as in (13). There is no Einstein-Hilbert term and there is thus no gravitation at all. It is also clear that this curvature-dependent term breaks gauge symmetries explicitly along with the loop-induced mass $c_V\Lambda_\wp^2$ [11, 10]. In view of these features, it is necessary to first determine how the GR emerges from within the effective action (12). Besides, it is necessary to also determine if emergence of the GR can kill the gauge symmetry breaking terms. These two points will be clarified below in the framework of affine gravity with distant inspiration from the Higgs mechanism.

Journal of Physics: Conference Series 2533 (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

4. Reminiscencing the Higgs Field: Affine Curvature

It is necessary to distinguish between two cases: A vector boson of mass M_V and a gauge boson of anomalous loop-induced mass (proportional to $\Lambda_\wp$). Firstly, mass M_V of a vector boson respects the Poincare symmetry simply because it is a Casimir invariant of the Poincare group [1, 9]. The loop-induced mass of a gauge boson, on the other hand, is a Poincare-breaking parameter since the cutoff $\Lambda_\wp$ itself breaks the Poincare symmetry [3]. Secondly, in the case of a massive vector boson, the goal is to restructure the vector field as a gauge field. In the case of a gauge boson with loop-induced mass, however, the goal is to kill the anomalous mass. Thirdly, being a Poincare-invariant quantity, M_V can be promoted to a suitable Higgs scalar Φ_V via

$$\Lambda_\wp^2 \mathrm{Tr}\left[V_\mu V^\mu\right] \longmapsto \Phi_V^\dagger V_\mu V^\mu \Phi_V \tag{14}$$

as in the Higgs mechanism [8, 9]. For a loop-induced gauge boson mass, however, scalar fields like the Higgs field cannot serve the purpose [8, 9] since the requisite field must be a Poincare-breaking one. In search for this Poincare-breaking field, one notes that, in a general second-quantized quantum field theory with no presumed symmetries, the Poincare (translation) invariance emerges if the Poincare-breaking terms are identified with the spacetime curvature [25]. Physically, this means that the Poincare-breaking sources in a QFT are the spots where curvature can emerge. It is all clear that the UV cutoff $\Lambda_\wp$ is the said Poincare-breaking source in an effective QFT. There is thus every reason to conclude that the sought-for Poincare-breaking field should be the spacetime curvature itself. In this regard, one notices that, in the effective action in (13), promotion of $\Lambda_\wp^2 g_{\mu\nu}$ to the Ricci curvature $R_{\mu\nu}({}^g\Gamma)$ as a Poincare-breaking spurion would completely eradicate the gauge symmetry breaking term $c_V \mathrm{tr}\left[V^\mu\left(\Lambda_\wp^2 g_{\mu\nu} - R_{\mu\nu}({}^g\Gamma)\right)V^\nu\right]$. But this promotion is inherently inconsistent simply because $R_{\mu\nu}({}^g\Gamma)$ does vanish in the flat spacetime limit while $\Lambda_\wp^2 g_{\mu\nu}$ does not. One way to resolve this inconsistency is to promote $\Lambda_\wp^2 g_{\mu\nu}$ to a curvature tensor that does not vanish in the flat spacetime limit. The simplest curvature as such is the affine Ricci curvature [13, 14, 15]

$$\mathbb{R}_{\mu\nu}(\Gamma) = \partial_\lambda \Gamma^\lambda_{\mu\nu} - \partial_\nu \Gamma^\lambda_{\mu\lambda} + \Gamma^\rho_{\lambda\rho}\Gamma^\lambda_{\mu\nu} - \Gamma^\lambda_{\rho\nu}\Gamma^\rho_{\mu\lambda} \tag{15}$$

of an affine connection $\Gamma^\lambda_{\mu\nu}$, which is a general connection independent of the metric tensor $g_{\mu\nu}$ and its Levi-Civita connection ${}^g\Gamma^\lambda_{\mu\nu}$. Needless to say, $\Gamma^\lambda_{\mu\nu}$ is not a tensor in curved spacetime but it transforms as a tensor in the flat spacetime in which coordinate transformations are the linear Lorentz transformations [26]. Now, having the affine curvature $\mathbb{R}_{\mu\nu}(\Gamma)$ at hand, the UV cutoff $\Lambda_\wp$ can be promoted to the affine curvature as

$$\Lambda_\wp^2 g_{\mu\nu} \longmapsto \mathbb{R}_{\mu\nu}(\Gamma) \tag{16}$$

in the same philosophy as the equation (14) in which $\Lambda_\wp$ is promoted to the Higgs scalar Φ_V. In essence, equation (16) is a map from UV cutoff to affine curvature, and possesses three key properties:

(i) It reduces to $\Lambda_\wp^2 \eta_{\mu\nu} \mapsto \mathbb{R}_{\mu\nu}(\Gamma)$ in the flat spacetime in which the affine connection acts as a tensor field [26],

(ii) It is meant to hold at the fundamental scale of gravity $M_{Pl} \equiv (8\pi G_N)^{-1/2}$, where G_N is Newton's gravitational constant.

(iii) It rests on the fact that both the affine curvature $\mathbb{R}_{\mu\nu}(\Gamma)$ and metric $g_{\mu\nu}$ are classical fields since effective QFTs like (12) cannot be extended with new quantum fields.

Then, the cutoff-to-curvature map (16), with these three key properties, takes the power-law correction action (13) to a metric-Palatini (metric-affine) action [15, 16]

$$\delta\overline{S}_{pow}[g,\digamma;\mathbb{R},\log\mu,R] = \int d^4x\sqrt{-g}\Bigg\{-\frac{M_{Pl}^2}{2}\mathbb{R}(g,\Gamma) - \frac{\overline{c}_O}{16}\left(\mathbb{R}(g,\Gamma)\right)^2 - \frac{c_\phi}{4}\phi^\dagger\phi\,\mathbb{R}(g,\Gamma) + c_V\mathrm{tr}\left[V^\mu(\mathbb{R}_{\mu\nu}(\Gamma) - R_{\mu\nu}({}^g\Gamma))V^\nu\right]\Bigg\} \quad (17)$$

in which $\mathbb{R}(g,\Gamma) \equiv g^{\mu\nu}\mathbb{R}_{\mu\nu}(\Gamma)$ is the affine curvature scalar. The parameters in this action deserve a detailed discussion: First, as follows from the equation (4), the fundamental scale of gravity (the Planck scale) can be defined as

$$M_{Pl}^2 = \frac{1}{2}\mathscr{M}^2(\mu = M_{Pl}) \xrightarrow{\text{one loop}} -\frac{1}{128\pi^2}\mathrm{str}\left[\overline{M}^2\log\left(\frac{\overline{M}^2}{M_{Pl}^2}\right)\right] \quad (18)$$

in which $\overline{M}^2 = M^2\big(\mu = M_{Pl}\big)$. This definition comes to mean that the gravitational scale is equal to $\mathscr{M}^2/2$ evaluated at the gravitational scale simply because the map (16) holds at the gravitational scale. It is clear that M_{Pl} remains put at its value in (18) since curvature is classical and matter loops have already been used up in forming the flat spacetime effective action in (1). It is also clear that gravity remains attractive if the bosonic sector is heavier ($\mathrm{str}\big[\overline{M}^2\big] > 0$) and if all the matter fields weigh below the gravitational scale ($\mathrm{str}\big[\overline{M}^2\big] \lesssim M_{Pl}^2$). In general, scalars, singlet fermions and vector-like fermions can weigh heavy without breaking gauge symmetries, and such heavy fields can dominate M_{Pl} through $\mathrm{str}\big[\overline{M}^2\big]$.

In line with the definition of the gravitational scale in (18), quadratic curvature coefficient $\overline{c}_O$ in (17) can be defined as

$$\overline{c}_O = c_O\left(\mu = M_{Pl}\right) \xrightarrow{\text{one loop}} \frac{(n_b - n_f)}{64\pi^2} \quad (19)$$

and this definition means that $n_b(n_f)$ is the number of bosons (fermions) having masses from zero way up to the gravitational scale. In other words, $n_b(n_f)$ comprises entirety of the bosons (fermions) since particles heavier than M_{Pl} are disfavored by the attractive nature of gravity. It is clear that $\overline{c}_O$ remains set at its value in (19) since curvature is classical and matter loops have already been used up in forming the flat spacetime effective action in (1).

5. Affine Dynamics: From UV Cutoff to IR Curvature

Having elucidated the parameters of the metric-Palatini action (17), it is time to go back to the total effective action (12). In fact, it takes the complete metric-Palatini form [16]

$$S_{eff}[g,\digamma;\mathbb{R},\log\mu,R] = S_{tree}[g,\digamma] + \delta S_{log}[g,\log\mu,\digamma] + \delta\overline{S}_{pow}[g,\digamma;\mathbb{R},\log\mu,R] \quad (20)$$

after using the action (17). This metric-Palatini theory contains both the metrical curvature $R({}^g\Gamma)$ and the affine curvature $\mathbb{R}(\Gamma)$. It remains stationary against variations in the affine connection (namely $\delta_\Gamma S_{eff}[g,\digamma;\mathbb{R},\log\mu,R] = 0$) provided that

$${}^\Gamma\nabla_\lambda\mathbb{D}_{\mu\nu} = 0 \quad (21)$$

such that ${}^\Gamma\nabla_\lambda$ is the covariant derivative of the affine connection $\Gamma^\lambda_{\mu\nu}$, and

$$\mathbb{D}_{\mu\nu} = \left(\frac{1}{16\pi G_N} + \frac{c_\phi}{4}\phi^\dagger\phi + \frac{\overline{c}_O}{8}g^{\alpha\beta}\mathbb{R}_{\alpha\beta}(\Gamma)\right)g_{\mu\nu} - c_V\mathrm{tr}\left[V_\mu V_\nu\right] \quad (22)$$

Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

is the disformal metric of tensor fields, including the affine curvature $\mathbb{R}(\Gamma)$. The motion equation (21) implies that $\mathbb{D}_{\mu\nu}$ is covariantly-constant with respect to $\Gamma^{\lambda}_{\mu\nu}$, and this constancy leads to the exact solution

$$\begin{aligned} \Gamma^{\lambda}_{\mu\nu} &= \frac{1}{2}(\mathbb{D}^{-1})^{\lambda\rho}\left(\partial_\mu \mathbb{D}_{\nu\rho} + \partial_\nu \mathbb{D}_{\rho\mu} - \partial_\rho \mathbb{D}_{\mu\nu}\right) \\ &= {}^{g}\Gamma^{\lambda}_{\mu\nu} + \frac{1}{2}(\mathbb{D}^{-1})^{\lambda\rho}\left(\nabla_\mu \mathbb{D}_{\nu\rho} + \nabla_\nu \mathbb{D}_{\rho\mu} - \nabla_\rho \mathbb{D}_{\mu\nu}\right) \end{aligned} \tag{23}$$

in which, needless to say, ${}^{g}\Gamma^{\lambda}_{\mu\nu}$ is the Levi-Civita connection of the curved metric $g_{\mu\nu}$. The Planck scale in (18) is the largest scale and therefore it is legitimate to make the expansions

$$\Gamma^{\lambda}_{\mu\nu} = {}^{g}\Gamma^{\lambda}_{\mu\nu} + \frac{1}{M_{Pl}^2}\left(\nabla_\mu \mathbb{D}^{\lambda}_{\nu} + \nabla_\nu \mathbb{D}^{\lambda}_{\mu} - \nabla^{\lambda}\mathbb{D}_{\mu\nu}\right) + \mathcal{O}\left(M_{Pl}^{-4}\right) \tag{24}$$

and

$$\mathbb{R}_{\mu\nu}(\Gamma) = R_{\mu\nu}({}^{g}\Gamma) + \frac{1}{M_{Pl}^2}\left(\nabla^{\alpha}\nabla_\mu \mathbb{D}_{\alpha\nu} + \nabla^{\alpha}\nabla_\nu \mathbb{D}_{\alpha\mu} - \Box \mathbb{D}_{\mu\nu} - \nabla_\mu\nabla_\nu \mathbb{D}^{\alpha}_{\alpha}\right) + \mathcal{O}\left(M_{Pl}^{-4}\right) \tag{25}$$

so that both $\Gamma^{\lambda}_{\mu\nu}$ and $\mathbb{R}_{\mu\nu}(\Gamma)$ contain pure derivative terms at the next-to-leading $\mathcal{O}\left(M_{Pl}^{-2}\right)$ order [11, 12]. The expansion in (24) ensures that the affine connection $\Gamma^{\lambda}_{\mu\nu}$ is solved algebraically order by order in $1/M_{Pl}^2$ despite the fact that its motion equation (21) involves its own curvature $\mathbb{R}_{\mu\nu}(\Gamma)$ through $\mathbb{D}_{\mu\nu}$ [13, 14]. The expansion (25), on the other hand, ensures that the affine curvature $\mathbb{R}_{\mu\nu}(\Gamma)$ is equal to the metrical curvature $R_{\mu\nu}({}^{g}\Gamma)$ up to a doubly-Planck suppressed remainder. In essence, what happened is that the affine dynamics took the affine curvature $\mathbb{R}$ from its UV value $\Lambda_{\wp}^2$ in (16) to its IR value R in (25). Indeed, in the sense of holography [27, 28], the metrical curvature R sets the IR scale [29] above which QFTs hold as flat spacetime constructs [21, 22, 23].

6. Restoration of Gauge Symmetries

One consequence of the solution of the affine curvature in (25) is that the problematic loop-induced gauge boson mass term gets defused as

$$\int d^4x\sqrt{-g}c_V \,\mathrm{tr}\left[V^{\mu}(\mathbb{R}_{\mu\nu}(\Gamma) - R_{\mu\nu}({}^{g}\Gamma))V^{\nu}\right] \xrightarrow{\text{equation (25)}} \int d^4x\sqrt{-g}\left\{\text{zero} + \mathcal{O}\big(M_{Pl}^{-2}\big)\right\} \tag{26}$$

after using the solution of the affine curvature in (25) in the metric-Palatini action (17) [10, 11, 12]. The $\mathcal{O}\big(M_{Pl}^{-2}\big)$ remainder here, containing the next-to-leading order derivative term in (25), involves derivatives of the scalars ϕ and gauge fields V_μ, and can produce therefore no mass terms for either of them. It is worth noting that the gauge symmetries broken explicitly by a Poincare-conserving (Poincare-breaking) UV cutoff are restored via the Higgs field Φ_V (via the affine curvature $\mathbb{R}$) at the minimum of the Φ_V potential energy (at the extremum of the metric-affine action). This contrast shows that Poincare-conserving and Poincare-breaking UV cutoffs are fundamentally different and lead, respectively, to field-theoretic and gravitational completions of the effective QFT.

7. Emergence of General Relativity

One other consequence of the solution of the affine curvature in (25) is that non-gauge sector of the metric-Palatini action (17) reduces to the quadratic curvature gravity

$$\begin{aligned} \int d^4x\sqrt{-g}\Bigg\{ & -\frac{M_{Pl}^2}{2}\mathbb{R}(g,\Gamma) - \frac{\overline{c}_O}{16}\left(\mathbb{R}(g,\Gamma)\right)^2 - \frac{c_\phi}{4}\phi^{\dagger}\phi\,\mathbb{R}(g,\Gamma)\Bigg\} \\ & \xrightarrow{\text{equation (25)}} \int d^4x\sqrt{-g}\left\{-\frac{M_{Pl}^2}{2}R - \frac{\overline{c}_O}{16}R^2 - \frac{c_\phi}{4}\phi^{\dagger}\phi\,R + \mathcal{O}\big(M_{Pl}^{-2}\big)\right\} \end{aligned} \tag{27}$$

Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

in which $R = g^{\mu\nu} R_{\mu\nu}({}^{g}\Gamma)$ is the usual curvature scalar in the GR. As in (26), the $\mathcal{O}\left(M_{Pl}^{-2}\right)$ remainder here consists of the next-to-leading order and higher terms in (25). It involves derivatives of the long-wavelength fields ϕ and V_μ, produces thus no mass terms for these fields, and remains small for all practical purposes.

The reductions (26) and (27) give rise to the total QFT plus GR action

$$S_{tot}[g, \digamma] = S_{tree}[g, \digamma] + \delta S_{log}[g, \log\mu, \digamma] + \int d^4x\sqrt{-g}\left\{-\frac{M_{Pl}^2}{2}R - \frac{c_O}{16}R^2 - \frac{c_\phi}{4}\phi^\dagger\phi R\right\} \quad (28)$$

in which the QFT sector

$$S_{\rm QFT}[g, \digamma] = S_{tree}[g, \digamma] + \delta S_{log}[g, \log\mu, \digamma] \quad (29)$$

is the usual $\overline{MS}$-renormalized QFT resting on the matter loops in flat spacetime and evolving from scale to scale by renormalization group equations in $\log\mu$ [3]. Its gravity sector

$$S_{\rm GR}[g, \phi] = \int d^4x\sqrt{-g}\left\{-\frac{M_{Pl}^2}{2}R - \frac{c_O}{16}R^2 - \frac{c_\phi}{4}\phi^\dagger\phi R\right\} \quad (30)$$

rests on the flat spacetime loop factors and emerges from the requirement of restoring gauge symmetries. With these two unique features, it differs from all the other matter-induced gravity theories (induced [30, 31, 32], emergent [33, 34], analogue [35], broken symmetry [36, 37], and the like). It is an $R + R\phi^2 + R^2$ gravity theory [38] whose each and every coupling is a flat spacetime loop factor (coefficient of $\Lambda_\wp^2$ or $\Lambda_\wp^4$ in (2)). It is the gauge symmetry-restoring emergent gravity or briefly the *symmergent gravity* [12, 11, 10], which is reformulated in a completely new setting in the present work. It should not be confused with the effective action computed in curved spacetime, which gives $M_{Pl}^2 \propto \Lambda_\wp^2$ along with $\Lambda_\wp$-sized scalar and gauge boson masses and a $\Lambda_\wp^4$-sized vacuum energy density [30, 31, 32]. Symmergent gravity, as reformulated and elucidated in the present work, stands out as a novel framework for completing effective QFTs in the UV when the UV cutoff is a Poincare-breaking one.

8. Discussions and Conclusion

In this talk, we have discussed how gauge symmetries broken explicitly by a Poincare-breaking UV cutoff can be restored, and what consequences such restoration can have. The end result is symmergent gravity as an emergent gravity plus $\overline{MS}$–renormalized QFT framework. Below, we discuss briefly some salient implications of the symmergent gravity [10, 11, 12].

First, **gravitational scale necessitates new physics beyond SM (BSM).** The gravitational scale M_{Pl} in $S_{\rm GR}(g, \phi)$ is induced by the flat spacetime matter loops as in (18). In the SM, it takes the value $M_{Pl}^2 \approx -G_F^{-1}$, where $G_F \approx (293 \text{ GeV})^{-2}$ is the Fermi scale. Its negative sign, set by the top quark contribution, is obviously unacceptable. It must be turned to positive if gravity is to be attractive, and this can be done if there exist new particles beyond the SM spectrum. These beyond-the-SM (BSM) particles are a necessity. In regard to the enormity of the gravitational scale (which can be brought to its physical shell by a conformal rescaling of the metric), the BSM sector must have

(i) either a light spectrum with numerous more bosons than fermions (for instance, $m_b \sim m_f \sim G_F^{-1/2}$ with $n_b - n_f \sim 10^{32}$),

(ii) or a heavy spectrum with few more bosons than fermions (for instance, $m_b \sim m_f \lesssim M_{Pl}$ with $n_b - n_f \gtrsim 10$),

(iii) or a sparse spectrum with net boson dominance.

Journal of Physics: Conference Series **2533** (2023) 012028 doi:10.1088/1742-6596/2533/1/012028

In general, BSM particles do not have to couple to the SM particles simply because all they are requiered to do is to saturate the super-trace in (18) at a value $M_{Pl}^2 \simeq \overline{M}_{Pl}^2$. In other words, there are no symmetries or selection rules requiring the SM particles to couple to the BSM particles. They can form therefore a fully-decoupled black sector [39, 40, 41, 11] or a feebly-coupled dark sector [11, 42, 43], with distinctive signatures at collider searches [44], dark matter searches [45], and other possible phenomena [11].

Second, **Higgs-curvature coupling can probe the BSM.** The loop factor c_ϕ in $S_{\rm GR}(g,\phi)$ couples the scalar curvature $R(g)$ to scalar fields ϕ. It is about 1.3% in the SM [10, 11]. Its deviation from this SM value indicates existence of new particles which couple to the SM Higgs boson. These BSM particles can be probed via their effects on various gravitational and astrophysical phenomena [46, 47, 48].

Third, **symmetries of the BSM sector might shed new light on the cosmological constant problem.** The vacuum energy contained in $S_{\rm QFT}(g,\digamma)$ [17]

$$V(\mu) \xrightarrow{\text{one loop}} V(\langle\phi\rangle) + \frac{1}{32\pi^2}{\rm str}\left[\overline{M}^4\left(1-\frac{3}{2}\log\frac{\overline{M}^2}{\mu^2}\right)\right] \tag{31}$$

gathers together field-independent $\log\mu$ corrections in (29) in the minimum $\phi = \langle\phi\rangle$ of the scalar potential $V(\phi)$. Its empirical value is $V_{emp} = \left(2.57\times10^{-3}\ {\rm eV}\right)^4$ [49]. The cosmological constant problem is to shoot this specific value with the prediction in (31), and such a shooting is tantalizingly fine-tuned [50]. But, as a way out possible only in symmergence, it might be possible to achieve a resolution if the BSM fields enjoy appropriate symmetries and selection rules [43]. For instance, a supersymmetric BSM sector would help in eliminating $V(\mu)$ but its realization can require extra structures [51].

Fourth, **quadratic curvature term can probe the BSM.** The loop factor c_O in $S_{\rm GR}(g,\phi)$ is proportional to the boson-fermion number difference. It vanishes identically in a QFT with equal bosonic and fermionic degrees of freedom (as in the supersymmetric theories [52, 53]) and, as a result, the gravitational sector in (30) reduces to the GR. This normally is not possible since under general covariance all curvature invariants can contribute to the gravitational sector, and it simply is not possible to get the exact GR. But symmergence is able to generate the GR, the exact GR, when the SM+BSM involves equal bosonic and fermionic degrees of freedoms [11, 12].

The loop factor c_O, when nonzero, acts as probe of the BSM in strong-curvature media. It can probe the BSM sector in terms of $n_b - n_f$ via strong-curvature effects. One such effect is the Starobinsky inflation, and seems to require $n_b - n_f \approx 10^{13}$ [54]. One other effect concerns the black holes, which put limits [55, 56, 57] on c_O and the vacuum energy $V(\mu)$ through the EHT observations [58].

Fifth, **heavy BSM does not necessarily destabilize the light scalars.** Light scalars ϕ_L in $S_{\rm QFT}(g,\digamma)$ are oversensitive to heavy fields. Indeed, their masses m_{ϕ_L} are shifted by an amount

$$\delta m_{\phi_L}^2 = c_{\phi_L}\lambda_{\phi_L\digamma_H} m_{\digamma_H}^2 \log\frac{m_{\digamma_H}^2}{\mu^2} \tag{32}$$

if they couple with loop factor c_{ϕ_L} and coupling constant $\lambda_{\phi_L\digamma_H}$ to heavy fields $\digamma_H$ of masses $m_{\digamma_H} \gg m_{\phi_L}$. This mass correction reveals that heavier the $\digamma_H$ larger the shift in the ϕ_L mass and stronger the destabilization of the light scalar sector. This is the famous little hierarchy problem [59]. It is the reason the null LHC results [60, 61] have sidelined supersymmetry and other known completions.

Can symmergence have a different say on the little hierarchy problem? Can it provide a resolution? To answer these questions, one notices that the BSM sector is necessitated for generating the gravitational scale, and all that its formula in (18) involves is super-trace over mass-squareds of the SM+BSM particles. In other words, there is no symmetry principle or selection rule requiring the SM and BSM fields to interact. Indeed, the coupling $\lambda_{\phi_L \digamma_H}$ is not under any constraint since workings of symmergence do not depend on it. This means that symmergence allows $\lambda_{\phi_L \digamma_H}$ be small enough to keep $\delta m^2_{\phi_L}$ small enough. More precisely, it is possible ensure $\delta m^2_{\phi_L} \ll m^2_{\phi_L}$ if $\lambda_{\phi_L \digamma_H}$ obeys the bound

$$|\lambda_{\phi_L \digamma_H}| \lesssim \lambda_{SM} \frac{m^2_{\phi_L}}{m^2_{\digamma_H}} \tag{33}$$

in which $\lambda_{SM} \sim \mathcal{O}(1)$ is a typical SM coupling. This "small-coupling domain" is specific to symmergence. It is the domain in which the SM and BSM are sufficiently decoupled and the little hierarchy problem is naturally avoided.

Symmergence, as a framework intertwining emergent gravity and new physics sector, can have widespread effects in various phenomena. Its study in collider, astrophysical and cosmological setting can provide therefore further information.

Acknowledgements
The author is grateful to his collaborators A. Abdujabbarov, İ. Çimdiker, C. Karahan, R. Pantig, B. Puliçe, J. Rayimbaev, A. Övgün and O. Sargın for their contributions to various results presented in this talk. The author thanks to participants of DICE 2022 for enlightening conversations and discussions.

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