A note on classical and quantum unimodular gravity

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We discuss unimodular gravity at a classical level, and in terms of its extension into the UV through an appropriate path integral representation. Classically, unimodular gravity is simply a gauge fixed version of General Relativity (GR), and as such it yields identical dynamics and physical predictions. We clarify this and explain why there is no sense in which it can “bring a new perspective” to the cosmological constant problem. The quantum equivalence between unimodular gravity and GR is more of a subtle question, but we present an argument that suggests one can always maintain the equivalence up to arbitrarily high momenta. As a corollary to this, we argue that whenever inequivalence is seen at the quantum level, that just means we have defined two different quantum theories that happen to share a classical limit.

I. INTRODUCTION

When Einstein laid down the foundations for General Relativity [1], he remarked that the laws of gravity sometimes took on a simpler form in certain coordinate systems, and illustrated his point by choosing so-called unimodular coordinates, where

\[ \det g_{\mu\nu} = -1. \] (1)

Of course, this choice of coordinates yields the same predictions as any other in a diffeomorphism invariant theory - a seemingly obvious fact that is at the heart of the equivalence between classical General Relativity and so-called unimodular gravity. It is particularly instructive to quote Einstein’s words on the unimodular choice of coordinates from this seminal paper [1]. There, after Einstein noted that this particular choice of coordinates simplify the derivation of the general relativistic equations, he characteristically went on to remark that ¹

“It would be mistaken to believe that this choice implies any sort of departure from the fundamental principle of GR. We don’t ask what are the (gravitational) laws which are invariant under all transformations that preserve the value of the determinant being one, but, instead, we ask which are the covariant laws, which after we derive, we can simplify their expression through a particular choice of coordinates.”

Unimodular gravity is obtained from a restricted variation of the Einstein-Hilbert action, in which the unimodular condition (1) is imposed from the beginning. The resulting field equations correspond to the traceless Einstein equations, and can easily be shown to be equivalent to the full Einstein equations with a cosmological constant term, \( \Lambda \), entering as an integration constant. Thus the equivalence to classical GR is made manifest, and there can be no sense in which unimodular gravity can say anything more or less than GR about anything to do with classical gravity. This includes the cosmological constant problem which is sometimes used as motivation for studying unimodular gravity [2–4] (see also e.g. [5–7]). One of the purposes of this paper is to make this point abundantly clear in a self-contained presentation, in the hope of addressing certain misconceptions that continue to appear in the literature.

Beyond classical gravity, however, the equivalence between GR and unimodular gravity is more subtle, with little consensus. For example, in Ref. [8] it has been argued that in principle quantum effects can allow one to discriminate between the two theories, while in Ref. [9] it was claimed that the two theories are equivalent at the perturbative level for asymptotically flat space times, but inequivalence was found for semiclassical non–perturbative quantities around particular backgrounds. In the canonical approach to quantum gravity, it has been suggested that unimodular gravity can help address the problem of time [10, 11], although such a claim has been strongly refuted [12]. We will present an argument suggesting that both theories can be extended into the quantum realm in such a way as to preserve their equivalence.

Intuitively, it is straightforward to see how the equivalence can be preserved. In the path integral formalism one must always divide out the symmetry group of the theory. For GR, the symmetry group is the diffeomorphism group (Diff), \( \delta g_{\mu\nu} = \nabla_{\mu} \xi_{\nu} + \nabla_{\nu} \xi_{\mu} \), whereas for unimodular gravity the condition (1) breaks this down to transverse diffeomorphisms (TDiff), satisfying \( \nabla_{\mu} \xi_{\nu} = 0 \). One might imagine taking the path integral for GR and first dividing out the longitudinal diffeomorphisms satisfying \( \nabla_{\mu} \xi_{\nu} \neq 0 \), such as to give us the path integral for unimodular gravity. This is essentially the spirit behind the claims made in Ref. [9], and we are certainly sympathetic to their approach.

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¹ A free translation from the original German text: “Es wäre aber irrtümlich, zu glauben, daß dieser Schritt einen partiellen Verzicht auf das allgemeine Relativitätspostulat bedeutet. Wir fragen nicht: „Wie heißen die Naturgesetze, welche gegenüber allen Transformationen von der Determinante 1 kovariant sind?“ Sondern wir fragen: „Wie heißen die allgemein kovarianten Naturgesetze?“ Erst nachdem wir diese aufgestellt haben, vereinfachen wir ihren Ausdruck durch eine besondere Wahl des Bezugssystems.”
Alternatively, we can always break the quantum equivalence between GR and unimodular gravity by force by defining them to be different from the beginning. For example, one can write unimodular gravity as a manifestly Diff (rather than TDiff) invariant theory by introducing extra (Stückelberg) fields [12, 13]. If the extra fields are source-free the equivalence to GR remains (at least classically), but if they are sourced, it is broken. This means that the equivalence can always be broken at the quantum level by allowing the additional fields to exist as external legs in Feynman diagrams.

We structure the paper as follows: In section II we clarify some fundamental aspects of unimodular gravity at the classical level and explain why it does not “bring a new perspective” to the cosmological constant problem. We also present a number of alternative formulations for a covariant unimodular action, some of which have not appeared in the literature before (to our knowledge). In section III we present a schematic argument that suggests that unimodular gravity and GR can be extended into the quantum realm in such a way as to maintain their equivalence, provided we make certain assumptions. We conclude in section IV.

II. CLASSICAL UNIMODULAR GRAVITY

Unimodular gravity is obtained from the Einstein-Hilbert action under a restricted variation that preserves the metric determinant,

$$\delta \frac{\sqrt{-g}}{\delta g_{\mu\nu}} = 0.$$  \hfill (2)

where \( g \equiv \det g_{\mu\nu} \). To understand the implications of the unimodularity condition (2) for the gauge symmetries of the theory, recall that GR is invariant under diffeomorphism transformations, infinitesimally described by

$$\delta g_{\mu\nu}(x) = \nabla_\mu \xi_\nu + \nabla_\nu \xi_\mu.$$  \hfill (3)

If we think of the metric variation in (2) as an infinitesimal, gauge transformation, using (3) one can see that the gauge vectors \( \xi_\mu \) are forced to satisfy the following transversality condition

$$\frac{1}{2} \sqrt{-g} g^{\alpha\beta} \delta g_{\alpha\beta} = 0 = \nabla_\mu \xi^\mu = 0.$$  \hfill (4)

Therefore, under condition (2) the set of allowed diffeomorphism transformations of the action is now restricted to those that satisfy the transversality condition (4), and as a consequence, the set of full diffeomorphism symmetries (Diff) is restricted to the subset of the transverse ones (TDiff).

The restriction (2) is most conveniently imposed using a scalar Lagrange multiplier \( \lambda(x) \), so that the action is given by [3, 4]

$$S = \int d^4x \left[ \sqrt{-g} \frac{R[g]}{16\pi G} - \lambda(x) (\sqrt{-g} - \epsilon_0) \right] + S_m$$  \hfill (5)

where \( S_m \) denotes the effective action for the (quantum) matter fields coupled to the (classical) metric, and \( \epsilon_0 \) is a non-dynamical volume element that explicitly breaks Diff down to TDiff. The resulting field equations yield

$$G_{\mu\nu} = 8\pi G T_{\mu\nu} - \frac{\lambda(x)}{2} g_{\mu\nu},$$  \hfill (6)

$$\sqrt{-g} = \epsilon_0,$$  \hfill (7)

where \( T_{\mu\nu} = -\frac{2}{\sqrt{-g}} \frac{\delta S_m}{\delta g_{\mu\nu}} \) is the effective energy-momentum tensor describing the matter fields. Taking the trace of equation (6) yields \( \lambda(x) = \frac{1}{4}(R + 8\pi G T) \), and the traceless Einstein equations follow, as expected. Furthermore, if we assume that the effective matter action is invariant under Diff, we have energy-momentum conservation, and taking the divergence of equation (6) we find that \( \partial_\mu \lambda = 0 \). This fixes the Lagrange multiplier to be a constant \( \lambda_0 \), so that the dynamics is equivalent to that of General Relativity with a cosmological constant, \( \lambda_0/2 \).

$$G_{\mu\nu} = 8\pi G T_{\mu\nu} - \frac{\lambda_0}{2} g_{\mu\nu}.$$  \hfill (8)

Because the cosmological constant enters as an integration constant, rather than a parameter in the action, it is often said that this brings a new perspective to the cosmological constant problem. Such statements are wholly nugatory, and fail to appreciate the true nature of the problem. We shall now clarify this.

The cosmological constant problem is usually described as follows: why is the observed value of Einstein’s cosmological constant at least sixty orders of magnitude less than that expected from vacuum energy contributions? In quantum field theory, the vacuum is well known to carry a non-trivial energy density, and a standard calculation reveals this to be at least \( \rho_{\text{vac}} \gtrsim (\text{TeV})^4 \). In the absence of gravity, one can simply define this as the zero point energy, and then ignore it as it does not enter the dynamics. When gravity is turned on, a combination of general covariance and the equivalence principle require that this vacuum energy should, like any other form of energy, gravitate. However, unlike more familiar sources of matter such as dust or radiation, the energy density of the vacuum stays constant in time and does not dilute with the expansion of the universe. In GR, the vacuum energy is combined with the bare cosmological constant \( \Lambda \), so that it is the combination, \( \Lambda + \rho_{\text{vac}} \), that actually gravitates. This combination should not exceed the critical density of the universe today, \( \frac{\Lambda}{8\pi G} + \rho_{\text{vac}} \lesssim (\text{meV})^4 \), requiring \( \Lambda \) to be fine-tuned to at least sixty decimal places.

The cosmological constant problem, as described above, is somewhat wrongly stated. As discussed in some detail in [14], the issue is not so much one of fine-tuning, but of radia-
tive instability. In quantum field theory, one regularly cancels off divergences in physical parameters before fixing any finite remainder empirically using observation. Indeed from the Wilsonian Renormalisation Group (RG) we know that the UV sensitivity of relevant operators renders them incalculable – they should be measured instead. This is just renormalisation in action - you pick the finite part of your counterterms to fit the observation. The real concern is when this renormalisation procedure becomes unstable against changes in the effective description, e.g. against additional loop corrections, or against changing the renormalisation group scale in the exact calculation of the Wilsonian effective action. In other words, by adding, say, additional loops, do I need to drastically retune the finite part of the appropriate counterterm (in this case, the bare cosmological constant)? In the Standard Model of Particle Physics, one typically finds that each additional loop correction adjusts the vacuum energy density by an amount $\Delta \rho_{\text{vac}} \gtrsim (\text{TeV})^4$, requiring the bare value of the cosmological constant to be retuned to the same level of precision. Similarly, in the Wilsonian action, the vacuum energy jumps by an amount $\Delta \rho_{\text{vac}} \gtrsim m^4$ whenever the cut-off passes through a threshold of mass, $m$, at least up to the TeV scale (for a nice discussion, see also [15]). Again this requires the bare cosmological constant to be retuned to considerable precision.

Now, whether we are working with the bare parameter $\Lambda$ from GR, or the integration constant $\lambda_0$ from unimodular gravity, the essence of the cosmological constant problem remains the same. To see this, imagine we define the effective action for matter (e.g. by specifying the order in (matter) loops in a perturbative description, or else up to some cut-off in the exact Wilsonian description) and compute the vacuum energy accordingly. We then tune $\Lambda$ or $\lambda_0$ to some high degree of precision. But what happens when the effective description is altered slightly (e.g., by changing the loop order in the perturbative case, or the location of the cut-off beyond a new mass threshold in the non-perturbative case)? A new computation of the vacuum energy yields a completely new value, and we are required to readjust $\Lambda$ or $\lambda_0$. In other words, a choice of $\Lambda$ in GR, or $\lambda_0$ in unimodular gravity is unstable against changing the effective field theory description of matter. Unimodular gravity does not bring any new degree of freedom to do that. The point, of course, is that CMB temperature anisotropies are observable and therefore independent of the choice of gauge. Therefore, in GR, we are free to perform the relevant calculations in unimodular gauge (1), not setting the shift variable to zero, inevitably recovering the same results as in unimodular gravity. It would appear that the technical reason for the discrepancy in [16] is that the Bianchi identity has not been properly accounted for. Another way of thinking about this is the following. Since after the enforcement of the unimodularity condition in the derivation of the Einstein equations, one can formally recover the standard GR equations with a cosmological constant at the covariant level, after taking into account the Bianchi identities, the result of the linearisation of the covariant equations around any background should yield exactly the same results as in GR.

Let us conclude this section by presenting some alternative formulations of unimodular gravity, all of which are classically equivalent. The first of these involves restoring the full diffeomorphism invariance in the action (5) by means of a Stückelberg trick. To this end we introduce four Stückelberg fields $\phi^\alpha(x)$, as if we were performing a general coordinate transformation, and let $x^\alpha \rightarrow x^\alpha + \phi^\alpha(x)$. The gravitational part of the action becomes [12],

$$S_{\text{stuck}} = \int d^4x \left[ -\frac{g}{16\pi G} \lambda \left(\sqrt{-g} - \epsilon_0 |J^\alpha_\beta|\right) \right], \quad (9)$$

where we have defined the determinant of the Jacobian matrix $J^\alpha_\beta \equiv \frac{\partial x^\alpha}{\partial \tau^\beta(1)}$ as $|J^\alpha_\beta| = 4! \delta^{[\alpha}_\mu \delta_{\nu]}^{\beta} \delta_{\lambda}^{\gamma} J^\mu_\alpha J^\nu_\beta J^\kappa_\gamma J^\lambda_\delta$. This is now explicitly invariant under diffeomorphisms $x^\mu \rightarrow x'^\mu(x^\mu)$, provided the Stückelberg fields, $\phi^\alpha$, transform as scalars$^3$. Furthermore, if we note that $|J^\alpha_\beta| = \partial_\alpha \left[ 4! \delta^{[\alpha}_\mu \delta^{[\beta}_\nu \delta^{[\gamma}_\lambda \phi^\mu J^\nu_\beta J^\kappa_\gamma J^\lambda_\delta \right]$ we see that the Stückelberg action is a special case of the Henneaux-Teitelboim action [13],

$$S_{\text{HT}} = \int d^4x \left[ -\frac{g}{16\pi G} \lambda \left(\sqrt{-g} - \partial_\mu \tau^\mu\right) \right], \quad (10)$$

where $\tau^\mu$ is a vector density. This action can be further generalised to

$$S_{\text{genHT}} = \int d^4x \sqrt{-g} \left[ -\frac{g}{16\pi G} \lambda f \left(\frac{\partial_\mu \tau^\mu}{\sqrt{-g}}\right) - q \left(\frac{\partial_\mu \tau^\mu}{\sqrt{-g}}\right) \right] \quad (11)$$

where the functions $f$ and $q$ do not fall into any of the following categories: (i) $f$ has no real zeroes; (ii) the isolated zeroes of $f$ and $f'$ coincide; (iii) $f$ is identically zero and $q$ is linear. The classical dynamics for this generalised action remains equivalent to that of GR with the cosmological

\footnote{Actually, the action (9) is invariant under Diff as long as $\phi^\alpha(x) \rightarrow \Phi^\alpha(\phi(x))$, where $\frac{\partial \Phi^\alpha}{\partial \phi^\beta} = 1$.}
constant entering as an integration constant. To our knowledge this generalised form of the unimodular action has not appeared in the literature before.

Assuming matter only couples directly to the metric, the generalised action (11) gives rise to the following field equations

\[ G_{\mu\nu} = 8\pi G [T_{\mu\nu} + g_{\mu\nu} (\lambda V(\psi) + U(\psi))], \]
\[ f(\psi) = 0, \]
\[ \delta_\alpha (\lambda f'(\psi) + q'(\psi)) = 0, \]

where \( T_{\mu\nu} \) is the energy momentum tensor of matter, \( \psi = \frac{\partial^2 \gamma^2}{\sqrt{\gamma}} \), and we have defined \( V(\psi) \equiv f^\prime(\psi) - f(\psi), \) \( U(\psi) \equiv \psi q'(\psi) - g(\psi) \). Equations (13) and (14) are constraint equations for the fields \( \lambda \) and \( \psi \), yielding some constant solution for \( \psi = \psi_0 \), and a constant, but arbitrary, Lagrange multiplier, \( \lambda \). When plugged into Einstein’s equations (12) this gives a constant, but arbitrary cosmological constant type term \( \lambda V(\psi) + U(\psi) \) on the RHS. Thus we recover the field equations of GR with a cosmological constant, as anticipated.

III. QUANTUM UNIMODULAR GRAVITY

In this section, we will provide a (less than rigorous) argument that unimodular gravity and GR can be extended into the quantum realm in such a way as to preserve their equivalence. To this end we start by defining the generating functional,

\[ Z[J] = \int Dg_{\mu\nu} D\lambda D\tau^\mu e^{iS_{HT}[g,\tau,\lambda] + iS_{\text{ext}}[g,J]} \]

where \( S_{HT} \) denotes the Henneaux-Teitelboim action (10) and \( S_{\text{ext}} \) the coupling to external sources. The full action is taken to be invariant under Diffs, as is the functional measure. Crucially, we have assumed that it is only the metric that couples to external sources and not the vector density, \( \tau^\mu \) or the scalar, \( \lambda \). Furthermore, the Henneaux-Teitelboim action has been endowed with a boundary term \([9],\)

\[ \int_{\text{boundary}} d^3x \sqrt{-\gamma} \left[ \frac{1}{8\pi G} K - \mu_\lambda \tau^\mu \right] \]

where \( \gamma \equiv \det \gamma_{\mu\nu} \) with \( \gamma_{\mu\nu} \) the induced metric on the boundary, \( \mu_\lambda \) the outward normal, and \( K \equiv K^\mu_{\mu} \) is the trace of the extrinsic curvature \( K_{\mu\nu} \). After integration by parts it is easy to see that \( \tau^\mu \) reduces to a Lagrange multiplier whose purpose is merely to fix \( \delta_\mu \lambda = 0 \). For a suitably chosen measure, the functional integration over \( \tau^\mu \) should yield

\[ Z[J] = \int Dg_{\mu\nu} D\lambda \delta_\mu \lambda e^{iS_{HT}[g,\lambda] + iS_{\text{ext}}[g,J]} \]

where

\[ S_{HT}[g,\lambda] = \int d^4x \sqrt{-g} \left[ \frac{R}{16\pi G} - \lambda(x) \right] + \int_{\text{boundary}} d^3x \sqrt{-\gamma} \frac{1}{8\pi G} K \]

In [9], it is argued that a physical boundary condition would be to impose no variation of \( \lambda \) at the boundary. This, along with the delta function, allows us to completely do the functional integration over \( \lambda \), yielding

\[ Z[J] = \int Dg_{\mu\nu} e^{iS_{GR}[g,\lambda_0] + iS_{\text{ext}}[g,J]} \]

where

\[ S_{GR}[g,\lambda_0] = \int d^4x \sqrt{-g} \left[ \frac{R}{16\pi G} - \lambda_0 \right] + \int_{\text{boundary}} d^3x \sqrt{-\gamma} \frac{1}{8\pi G} K \]

and \( \lambda_0 \) is the arbitrary fixed boundary value of \( \lambda \). Thus we arrive at the generating functional for General Relativity with a cosmological constant \( \lambda_0/2 \). Our somewhat schematic argument strongly suggests that there is a clear way in which we can extend unimodular gravity in to the UV so that it maintains its equivalence to GR.

What if we do not fix \( \lambda \) on the boundary? Then the delta function in Eq. (17) does not allow us to completely do the functional integration over \( \lambda \). In particular we are left with an ordinary integration over space-time constants \( \lambda_0 \),

\[ Z[J] = \int Dg_{\mu\nu} d\lambda_0 f(\lambda_0) e^{iS_{GR}[g,\lambda_0] + iS_{\text{ext}}[g,J]} \]

where we have included a possible non-trivial contribution, \( f(\lambda_0) \), to the measure for completeness. Classically this suggests a theory which is locally equivalent to GR but with an additional global constraint coming from variation over a global parameter - the bare cosmological constant (in this case, \( \lambda_0/2 \)). This is somewhat reminiscent of (although not identical to) the recent scenario [14, 17] in which the Standard Model contribution to vacuum energy is successfully sequestered from gravity to all orders in Standard Model loops.

For the Diff invariant manifestations of unimodular gravity discussed above, the arguments in support of quantum equivalence to GR relied heavily on the assumption that the additional fields, \( \lambda(x) \), \( \tau^\mu(x) \) are source-free and can be eliminated as auxiliary fields. We expect to break this equiva-
lence the moment we switch on corresponding sources, or in other words, we allow $\lambda(x)$ and $\tau^\mu(x)$ to lie on the external legs of Feynman diagrams. However, we emphasise that by doing this we are breaking the quantum equivalence by hand.

IV. DISCUSSION AND CONCLUSIONS

There are three messages we would like the reader to take away from this paper:

1. Classical unimodular gravity = classical GR, so any suggestion that the former can shed new light on any problems faced by the latter are entirely nugatory.

2. Quantum unimodular gravity = quantum GR provided we make certain assumptions about how we extend into the UV.

3. Quantum unimodular gravity $\neq$ quantum GR if we break those assumptions, but that is our choice, and amounts to defining the theories to differ in the UV.

The classical equivalence and its implications are spectacularly obvious, but confusion continues to reign in the literature. Perhaps this would not have happened if instead of being unimodular gravity, the theory were called gauge fixed GR. Of course, gauge invariance is merely a redundancy of description, and when the (residual) gauge invariance of unimodular gravity and GR are stripped away by a complete gauge fixing, it is inevitable that we always end up with the same physical predictions. In this paper we have discussed the implications of this in the context of unimodular gravity's favourite plaything, the cosmological constant problem, and we have explained in detail how it does not "bring a new perspective" to the problem in any way. Of course, this had to be true, because you do not change classical physics by choosing a gauge!

The quantum equivalence between the two theories is more of a subtle issue. We have presented a schematic argument based on the path integral approach to quantum gravity that suggests one can always maintain equivalence up to arbitrarily high momenta. The argument uses covariant descriptions of unimodular gravity [12, 13], where additional fields can be rendered purely auxiliary such that they may be integrated out in the path integral leaving us with the path integral for GR, with appropriate boundary conditions.

Whenever the quantum equivalence is seen to fail, we would argue that this says more about how one chose to go about extending the theories into the UV, than some inevitable inequivalence at the quantum level. Indeed, that choice amounts to defining the quantum theories to be inequivalent. Our discussion here has focussed on the path integral approach to quantum gravity, but we would not be at all surprised if this philosophy also extended to other approaches, although a precise demonstration of this may be less straightforward.

To sum up then: classical unimodular gravity and classical GR are the same thing, and they can be extended into the UV such that the equivalence is maintained. Whenever inequivalence is seen at the quantum level, that just we means we have defined two different quantum theories that happen to share a classical limit. An example of the latter is given by [12] with the Stückelberg fields and Lagrange multiplier allowed to lie on external legs.

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