The substantial role of Weyl symmetry in deriving general relativity from string theory

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Abstract

String theory reduces to general relativity in appropriate regimes. Huggett and Vistarini have given an account of this reduction that includes a deflationary thesis about symmetry: though the usual derivation of general relativity from string theory appeals to a premise about the theory's symmetry, Huggett and Vistarini argue that this premise plays no logical role. In this paper I disagree with their deflationary thesis and argue that their analysis is based on a popular but flawed conception of the interaction between symmetry and quantization. On this conception, quantization can break symmetries of the classical theory, and we must decide whether these symmetries should be reimposed. I argue that a better conception recognizes a tripartite distinction between ordinary, broken, and anomalous symmetries.

1 Introduction

The basic conceptual task for a quantum theory of gravitation is to recover something like the rough-and-ready picture of space and time that we use to characterize the target gravitational phenomena. Much recent philosophical work on this topic addresses general questions about this task: the extent to which theories must presuppose the rough-and-ready picture, the different possible success conditions for a recovery, whether and how to make sense of claims that spacetime emerges from some more fundamental quantum features of the world, and so on (Huggett and Wüthrich, 2013; Crowther, 2018). But there are also specific questions that arise within particular research programs. For example, Huggett and Vistarini (2015) and Vistarini (2019) point out that symmetry considerations seem to be "a key concept connecting string theory to phenomenological space-time" (2015, 1170) but that the status of these considerations is obscure. Despite the apparent importance of symmetry, Huggett and Vistarini argue that it is merely a formal feature of string theory, suggesting that it can play no substantial role. Getting to philosophical grips with string theory as a quantum theory of gravity requires a resolution of this conceptual tension.

This paper disagrees with Huggett and Vistarini's account and suggests an alternative. Their argument focuses on the interaction between symmetry and quantization. We recover the Einstein field equation (EFE) in string theory by quantizing a classical theory that exhibits so-called Weyl symmetry. On Huggett and Vistarini's telling, quantization breaks this Weyl symmetry, and we face a choice: if we decide to reimpose the symmetry then we are led to the EFE, and if we try to leave the symmetry broken then it will reappear in a different guise and again lead to the EFE. They conclude that Weyl symmetry is unavoidable and hence "not a logically independent postulate" (2015, 1173) of string theory. But I think this framing is misleading. When a theory is quantized, its symmetries have three possible fates: they might be preserved, they might be broken, or they might be anomalous. As I will argue, Huggett and Vistarini confine their attention to the first case; that is, their conclusion that Weyl symmetry is always preserved rests on considering only those cases in which it is preserved. Further attention to the other cases—and especially to cases in which the symmetry is anomalous—shows that Weyl symmetry is not a merely formal feature of string theory. It also illustrates the more general use of symmetries in the string-theoretic approach to theory construction, which plays an important role in defenses of the string theory program (Dawid, 2013).

The plan is as follows. In Section 2 I isolate the feature of Huggett and Vistarini's framing that I disagree with. Though this debate is motivated by the string-theoretic derivation of the EFE, my disagreement with Huggett and Vistarini is really a disagreement about the interpretation of quantum field theories with spacetime-dependent symmetries. It bears on their analysis only insofar as the derivation they discuss occurs within such a theory. As I argue, their presentation supposes that classical symmetries are either broken or preserved in the process of quantization. Section 3 argues that this is not a natural dichotomy, for there is a third possibility: the symmetry might be anomalous. Anomalous symmetries are in some sense preserved and in some sense broken; as such, they fit uncomfortably in Huggett and Vistarini's account. Section 4 uses this third category to argue that Weyl symmetry is not merely a formal feature of string theory.

2 Weyl symmetry

The argument that string theory reproduces general relativity in the appropriate domains has various prongs; Weyl symmetry is primarily relevant to the recovery of the EFE for the spacetime metric. In the appropriate regimes, a morass of string excitations ought to look like a Lorentzian metric to a test string moving through it. If the test string is to have a consistent quantization it must be Weyl invariant, and if it is to be Weyl invariant then the effective Lorentzian metric must satisfy the EFE. Or at least, this is the standard story. Huggett and Vistarini ultimately argue that this appeal to Weyl invariance can be circumvented: whether or not we suppose Weyl invariance, the EFE will follow. This is the claim I want to take issue with.

On Huggett and Vistarini's analysis, string theory's recovery of general relativity—and, thereby, phenomenological space—is expressed by two results. First, the spectrum of the string contains gravitons, the force carriers for gravity. More precisely, upon quantizing the string you will find quantum states containing massless spin-2 particles, the representation of the Lorentz group in which gravitons live. Second, an adequate quantization of a string moving in an approximately classical background made up of these massless spin-2 quanta requires the background to satisfy the EFE. String theory therefore contains the right stuff behaving in the right way to reproduce general relativity in the right regimes.

This paper is concerned with the second of these results, which is set within an effective theory of a string propagating in a classical background. Effective field theories model the salient degrees of freedom in systems where the fundamental degrees of freedom are unknown, or cannot be connected to the salient degrees of freedom, or are computationally intractable. For example, the Standard Model of particle physics contains twelve elementary matter particles. While some of these particles are observable in isolation at low energies, some only appear in bound states—the up and down quarks only occur as constituents of protons, neutrons, and pions. The effective degrees of freedom at low energies are therefore not those appearing in the Standard Model, and the effective field theory used to model physics at this scale is formulated directly in terms of protons, neutrons, and pions instead of quarks. Analogously, the EFE is derived from string theory in an effective theory that replaces gravitonic excitations with a classical Lorentzian metric. It's this metric that must satisfy the EFE.

More formally, the effective theory of interest is the following. The classical background is given by a metric G on a manifold X of dimension D. A possible history for a string is a map $\Sigma \to X$ with Σ a two-dimensional surface. We define a quantum field theory on Σ with two fluctuating fields: a Lorentzian metric g on Σ and a map $\phi : \Sigma \to X$ picking out a possible history for the string. The action for this theory is

$$S(g, \phi) = \int_\Sigma \|d\phi\|^2 \, \mathrm{vol}_g$$

where vol_g is the volume element on Σ determined by g and the norm is induced by g and G. That is, picking some coordinates on Σ and X, we have

$$\|d\phi\|^2 = g^{mn} \left(\partial_m \phi^\mu\right) \left(\partial_n \phi^\nu\right) G_{\mu\nu} \qquad\qquad \mathrm{vol}_g = \sqrt{-|g|}\, d^2x$$

where Roman indices run over the two dimensions of Σ and Greek indices over the D dimensions of X.

The EFE for G is obtained by requiring the quantum theory to be well-behaved. As a first pass at articulating this requirement, consider the path integral quantization of the action above. The theory is determined by the path integral

$$\int \mathcal{D}g \int \mathcal{D}\phi \, \exp(iS(g, \phi))$$

If we fix a metric g on the worldsheet Σ, the inner integral is the path integral for a quantum field theory of D scalar fields in two spacetime dimensions. Theories of this kind are relatively well understood, and the integral over ϕ is relatively easy to perform, at least when G is nearly flat. The path integral of our effective theory is therefore an integral over a family of scalar field theories indexed by metrics g on Σ. Integrating out the scalar fields, our integral becomes

$$\int \mathcal{D}g \, \exp(iS_{\text{red}}(g))$$

where $S_{\text{red}}(g)$ is an action for g that incorporates the quantum fluctuations of the field ϕ. The full path integral ought therefore reduce to an integral over g alone.

We can only integrate out the field ϕ if G satisfies the EFE. If the reduced action S_{red} exists then its exponentiation must have the same symmetries as the integral over ϕ. In particular, note that the original action $S(g, \phi)$ is invariant under the Weyl transformation

$$g \mapsto e^{2\omega} g$$

determined by a positive real-valued function ω on Σ, since the induced change in the norm of $d\phi$ cancels out the induced change in the volume element. On the other hand, to leading order in fluctuations in ϕ, an infinitesimal Weyl transformation with parameter ω shifts S_{red} by (D'Hoker, 1999, Eq. 6.61)

$$\delta S_{\text{red}}(g) = -\frac{1}{2\pi} \int_{\Sigma} d^2x \, \sqrt{-|g|} \, \omega \left(\frac{1}{2} g^{mn} (\partial_m \phi^\mu)(\partial_n \phi^\nu) R^G_{\mu\nu} + \frac{D-26}{6} R^g \right)$$

with $R^G_{\mu\nu}$ the Ricci tensor associated with the metric G and R^g the scalar curvature associated with the metric g. Since S_{red} must have the same symmetries as the original theory, this shift must vanish for all ω and g. This implies that the first term in the integrand vanishes for all g and ϕ, so that $R^G_{\mu\nu} = 0$. And this is the EFE in vacuum. The argument generalizes: if we add other classical background fields on X to the effective action then the shift in S_{red} under a Weyl transformation will include terms involving these other fields, and the shift will vanish when $R^G_{\mu\nu}$ satisfies the EFE determined by the stress-energy tensor of the added fields.

Huggett and Vistarini argue that Weyl symmetry plays no substantial role in the derivation I've just sketched; this is our point of disagreement. The derivation relied on the claim that S_{red} must have Weyl symmetry, and this "must" requires justification. Huggett and Vistarini argue that it is tautologous:

> although the derivation of the EFEs appeals to [Weyl] symmetry, since that is itself a consequence of string theory, it is not, logically speaking, a necessary premise of the derivation (2015, 1173).

On their view, the appeal to Weyl symmetry could in principle be eliminated. We must have $R^G_{\mu\nu} = 0$ (and $D = 26$), and S_{red} must be Weyl-invariant, but according to Huggett and Vistarini this is a downstream consequence of other hypotheses in string theory. We could just as well take a different route, one that made no explicit detour through Weyl symmetry.

Weyl symmetry certainly appears to play a role in the derivation just sketched, so Huggett and Vistarini argue for their triviality thesis by arguing that this is a mere appearance. In light of the generally nontrivial behavior of S_{red} under Weyl transformations, Huggett and Vistarini say that the Weyl symmetry of the original action is "broken by quantization" and that the EFE appears to follow when the symmetry is "reimposed" on S_{red} (2015, 1170). If we don't reimpose the symmetry then it seems we might have $R^G_{\mu\nu} \neq 0$ or $D \neq 26$. But, they claim, if we don't demand Weyl invariance then we must change our classical background:

> In this case, different choices of conformal factor in the Weyl transformation of the internal metric... will be physically different. Hence, $[\omega]$ is a new physical degree of freedom over the worldsheet, a scalar background field: specifically a dilaton field $[\Phi]$ (2015, 1171).

Suppose, then, that we adopt a different effective field theory, one that includes a scalar field Φ on X. Huggett and Vistarini argue that if we suppose Φ to be tachyonic, and if we suppose that some mechanism gives it good long-distance behavior, then we can show that Weyl invariance must hold. They conclude that Weyl invariance is unavoidable: even if we suppose that it doesn't hold we can derive that it does.

In the rest of this paper I argue that Huggett and Vistarini's reasoning does not go through and that talk of breaking and reimposing Weyl invariance is misleading. The Weyl symmetry of S_{red} is a necessary premise in the derivation of the EFE just sketched and is not a logical consequence of some other hypotheses. Huggett and Vistarini's argument only shows that S_{red} is Weyl-invariant under the hypothesis that S_{red} is Weyl-invariant. The triviality of this conclusion is obscured by a common way of talking about the role of symmetries in quantization, according to which quantization can break symmetries of the classical theory and it's left for us to decide whether to reimpose them. A better accounting of the situation distinguishes between cases in which the symmetry is preserved, cases in which it cannot be implemented, and cases in which it is anomalous.

3 Anomalies

My disagreement with Huggett and Vistarini's framing isn't particular to Weyl symmetry but applies to symmetries of all kinds. This section illustrates an alternative framing according to which any symmetry might be preserved, broken, or realized anomalously. Huggett and Vistarini also take their discussion to generalize to other kinds of symmetry. They explicitly analogize Weyl and gauge symmetry, and elsewhere Vistarini suggests that the possibility of a substantial role for Weyl symmetry "challenges the general idea that gauge symmetries are simply formal features of the way in which a theory's physical content is formally represented" (2019, 40).[1] I agree that deflationary views about Weyl and gauge symmetry stand and fall together, but I think they are untenable in both cases. They fail to mark an important distinction between anomalous and broken symmetries.[2]

3.1 Anomalous global symmetries in field theories

There's an important difference between a theory's being invariant under a symmetry or anomalous, and both of these situations are importantly different from a theory lacking that symmetry altogether. These differences can be illustrated by simpler theories with obvious physical application. Anomalous symmetries are also found in quantum field theory, where they can play an important role in saving the phenomena. The chiral anomaly in the Standard Model is a relatively simple example that illustrates why anomalous global symmetries are acceptable.

A particularly simple instance of the chiral anomaly appears in quantum electrodynamics with one charged fermion. The setting is four-dimensional Minkowski space, and the two fields in the theory are the electromagnetic gauge potential A and a massless Dirac fermion ψ. The action is

$$S(A,\psi) = \int_{\mathbb{M}^4} d^4x \left(-\frac{1}{4} F_{\mu\nu}F^{\mu\nu} - \bar{\psi} \not{D} \psi \right)$$

with $F_{\mu\nu}$ the field strength and $\not{D}$ the Dirac operator determined by A. As in the string theory of Section 2, the theory is specified by an iterated path integral

$$\int \mathcal{D}A \int \mathcal{D}\psi \, \mathcal{D}\overline{\psi} \, \exp(iS(A,\psi))$$

As before, we think of the inner integral as defining a quantum field theory with a single fluctuating fermion field ψ in the presence of a fixed classical electromagnetic potential A. And again as before, we proceed by integrating out the fermionic degrees of freedom to obtain an action depending only on A, with the full theory given by performing this remaining integration.

The integral over the fermion field transforms anomalously under the global symmetries of the action. Recall that a Dirac fermion ψ naturally decomposes into two parts: a left-handed Weyl fermion and a right-handed Weyl fermion. The Dirac operator $\not{D}$ is chirally symmetric, so the fermion term in $S(A,\psi)$ can be split into two terms, one involving the left-handed component of ψ and one involving the right-handed component. The action $S(A,\psi)$ is therefore invariant under two kinds of phase transformations,

$$\psi \mapsto e^{i\theta}\psi \qquad\qquad \psi \mapsto e^{i\theta\gamma^5}\psi$$

[1]See Healey (2007) and Redhead (2003) for more detailed articulations of this general idea as well as some discussion about how issues of symmetry and quantization are related to more obviously philosophical issues.

[2]What follows are two simple examples of anomalies. See Monnier (2019) for more thorough but still relatively informal discussions of anomalous quantum field theories.

the first of which rotates the phase on the left- and right-handed components by the same angle θ, and the second of which rotates the phases of each component the same magnitude θ but in opposite directions. Call the latter a chiral phase rotation, since it treats left- and right-handed components differently. While the integral over the fermion fields invariant under the first type of phase rotation, a chiral rotation by θ shifts the quantum effective action by

$$\delta S_{\text{eff}}(A,\psi) = -\frac{Q^2}{16\pi^2}\int_{\mathbb{M}^4} d^4x\,\theta\epsilon^{\mu\nu\alpha\beta}F_{\mu\nu}F_{\alpha\beta}$$

where Q is the charge of the fermion. The effective fermion action transforms under chiral rotations, so there's a sense in which it exhibits chiral symmetry. But it is not invariant; it has an anomaly.

The chiral anomaly doesn't vanish, but this isn't a problem. The derivation in Section 2 required the Weyl anomaly to vanish, but a vanishing chiral anomaly would lead to empirical inadequacy (Dougherty, 2020). For example, neutral pions decay to photons at a rate proportional to the chiral anomaly. If the chiral anomaly vanished then pions would hardly ever decay to two photons, but this is their most common decay channel. Indeed, the chiral anomaly was first discovered when trying to account for the neutral pion's decay rate. As another example, the mass of the η' meson is approximately proportional to the chiral anomaly. A theory without the chiral anomaly gets the η' meson's mass wrong by almost an order of magnitude.

The effective action varies under chiral rotations, but it exhibits chiral rotation symmetry in a weaker sense. Because the chiral anomaly is a reflection of this weaker invariance, it reflects the structure of the symmetry group by satisfying the so-called Wess–Zumino consistency conditions. This is importantly different from a theory that is not invariant under chiral rotations at all, like a theory with massive fermions. These two cases should be distinguished.

3.2 Anomalous gauge symmetries in field theories

Anomalous global symmetries like the chiral symmetry of Section 3.1 or Galilei symmetry in nonrelativistic quantum mechanics are unobjectionable. Indeed, they are desirable, because neutral pions often decay and the mass of a non-relativistic particle isn't state-dependent. Anomalous spacetime-dependent symmetries are less anodyne. These include Weyl symmetry when $R^G_{\mu\nu} \neq 0$ or $D \neq 26$, but they are also found in minor modifications of the Standard Model. The demand for a vanishing Weyl anomaly is analogous to the demand for a vanishing gauge anomaly, and the latter is perhaps more easily interpreted in physical terms by comparison with the Standard Model.

To illustrate gauge anomalies, consider a slightly different theory of charged matter. Replace the Dirac fermion in the action of Section 3.1 with a charged left-handed Weyl fermion χ to give the action $S(A,\chi)$. This action exhibits a spacetime-dependent symmetry: for any real-valued function α on $\mathbb{M}^4$ the transformation

$$A_\mu \mapsto A_\mu - \partial_\mu\alpha \qquad\qquad \chi \mapsto e^{iQ\alpha}\chi$$

leaves the action $S(A,\chi)$ unchanged. But when we integrate out the fermion χ the reduced action transforms anomalously:

$$\delta S_{\rm red}(A) = -\frac{Q^3}{96\pi^2}\int_{\mathbb{M}^4} d^4x\,\alpha\epsilon^{\mu\nu\alpha\beta}F_{\mu\nu}F_{\alpha\beta}$$

This resembles the chiral anomaly but is distinct. The action $S(A,\chi)$ doesn't have chiral symmetry at all—neither ordinary nor anomalous—because it contains only a left-handed Weyl fermion. And integrating out the Dirac fermion ψ from the action $S(A,\psi)$ of Section 3.1 produces a gauge-invariant reduced action, not one that transforms anomalously under gauge transformations.

The Standard Model has no gauge anomalies, and this expresses a nontrivial fact about the charges of its various particles. If we replaced the Weyl fermion in $S(A,\chi)$ with a right-handed fermion of the same charge then we would obtain another anomalous theory, and the anomaly would have the same form but with a sign flip. So we can build a theory with no anomaly if we include two Weyl fermions with the same charge, one of each handedness, for then the two anomaly terms would cancel. This is just the theory of Section 3.1, since a Dirac fermion is a pair of opposite-handed Weyl fermions. By similar reasoning, the anomaly associated with the $U(1)$ hypercharge gauge symmetry in the Standard Model is proportional to (Schwartz, 2014, Eq. 30.73)

$$2(Y_L^3 + 3Y_Q^3) - (Y_e^3 + Y_\nu^3 + 3(Y_u^3 - Y_d^3))$$

where the Ys are the hypercharges of left-handed leptons and quarks and right-handed electron, neutrino, and up- and down-type quarks. In the classical action these charges are freely specifiable independently, but their observed values are such that this expression vanishes. Similar anomaly cancellation conditions hold for other gauge symmetries in the Standard Model. And, of course, these are all cousins of the Weyl anomaly cancellation conditions $R^G_{\mu\nu} = 0$ and $D = 26$.

The examples in this section show that the preserved–broken dichotomy Huggett and Vistarini employ is too coarse a classification. Putting anomalous symmetries in the "broken" bucket neglects the fact that they satisfy nontrivial constraints, like the Wess–Zumino consistency condition. But putting them in the "preserved" bucket erases the difference between cases where anomalies cancel and cases where they don't. In particular, it elides theories where the Weyl anomaly vanishes and theories where it doesn't. Once we recognize that symmetries may be anomalously realized, we can further distinguish between anomalous global symmetries, like chiral symmetry, and anomalous gauge symmetries, like that of an electromagnetically charged Weyl fermion or the Weyl symmetry of Section 2's string theory. While the former obtain in perfectly good theories—both in principle and of particle phenomena—the latter are ruled out in the string-theoretic derivation of the EFE.

4 Theory space

The distinctions introduced in Section 3 clarify the task of justifying the string-theoretic derivation of the EFEs, and they lead to a problem for Huggett and Vistarini's analysis. The desired conclusion of $R^G_{\mu\nu} = 0$ follows from the demand that the Weyl anomaly vanish, and this follows from the demand that the total gauge anomaly always vanish. So we need a justification for this more general demand. I won't try to provide one here. But, supposing this demand is justified, it's a nontrivial one. There are theories that exhibit a nonvanishing gauge anomaly, like electrodynamics with a single Weyl fermion, and there are theories in which the gauge anomaly vanishes. Anomaly cancellation isn't tautologous. Indeed, the stringency of anomaly cancellation is sometimes claimed to uniquely determine a possible model of string theory.

Weyl symmetry plays a substantial role in Section 2's derivation of the EFE. This derivation required the reduced action to be exactly Weyl symmetric, not anomalously so. This is a substantial requirement because it forces us to have $R^G_{\mu\nu} = 0$ and $D = 26$. And this requirement is nontrivial because there are metrics that aren't Ricci flat and there are manifolds with dimension other than 26. In just the same way, demanding gauge anomaly cancellation in the theories of Section 3 or in the Standard Model puts nontrivial constraints on the field content and charges. It rules out a theory containing a single charged Weyl fermion, and it requires the electron's charge to be precisely the opposite of the proton's. Far from being a tautology, the vanishing of the Weyl anomaly is a powerful constraint on the construction of a quantum field theory.

The power of the vanishing anomaly condition requires an equally powerful justification, and I think this deserves further philosophical attention. Certainly we can't count every anomaly as pathological, since the chiral anomaly in Section 3.1 is instrumental in reproducing low-energy collider phenomena. But we can demand that every gauge anomaly vanish, and this demand is often made. It is sometimes said that theories with gauge anomalies aren't "coherent" (Dawid, 2013, 12) or "consistent" (Schwartz, 2014, 627), but this isn't obviously right, at least not in the strict sense. The traditional argument for this conclusion claims that gauge anomalies "destroy the renormalizability, and thus the consistency, of the gauge theory" (Bertlmann, 1996, 245). This seems too quick. Plenty of perfectly respectable theories aren't renormalizable, including the effective field theories used to model low-energy collider physics (Weinberg, 1995, §12.3). On the other hand, these effective field theories have unitary truncations at each order in the momenta, while any finite truncation of a gauge theory spoils unitarity. This is not the place to sort out the exact relationship between gauge anomalies and renormalization, but this relationship should be clarified if we would like to better understand the derivation of the EFE in string theory.

Because the vanishing Weyl anomaly is a nontrivial constraint, Huggett and Vistarini's deflationary argument must misfire. The problem with it is clear if we adapt it to a simpler theory with anomalous gauge symmetry, like the theory of the single charged fermion. Their argument, recall, begins by supposing that the reduced action isn't exactly invariant under the gauge symmetry. It's a matter of mathematical fact that the reduced action transforms under the gauge symmetry; the only question is whether it's invariant or anomalous. If we suppose it's anomalous then the fermion's charge must be nonzero. The theory then has a gauge anomaly and is not invariant under the gauge symmetry. At this point Huggett and Vistarini introduce new degrees of freedom and show that these cancel the anomaly. The analogous move in our charged fermion theory would be the introduction of further Weyl fermions: one fermion with the opposite handedness and the same charge, or two fermions with the opposite handedness and charge $Q/2$, or one Weyl fermion with the same charge and handedness and two with the same charge and opposite handedness, or something like this. The total gauge anomaly in any of these modified theories vanishes, so they are exactly gauge-invariant. But they're also just different theories. Introducing another fermion doesn't make the theory with one fermion consistent, it gives a theory with two fermions. In the same way, Weyl invariance in a theory containing a background scalar field Φ doesn't lead to Weyl invariance in a theory without a background scalar field.

Huggett and Vistarini's reasoning doesn't show that Weyl invariance is a purely formal requirement, but it can be useful in a different way. Anomaly cancellation can be a guide to theory development, because it can suggest modifications for the sake of anomaly cancellation. If you observe a charged Weyl fermion then there must be at least one more out there, because a theory with only one charged Weyl fermion has a gauge anomaly. Anomaly cancellation therefore constrains our exploration of the possible space of theories. Dawid's (2013) account of non-empirical theory assessment promotes this type of constraint to a general method for evaluating scientific theories. If we have reason to believe that the vast majority of theories have gauge anomalies then the fact that we've found some that lack them—the Standard Model, or the string theory with $R^G_{\mu\nu} = 0$ and $D = 26$—is a good sign that we're on the right track. The antecedent is a big "if", but it does seem difficult to construct theories in which all anomalies cancel.

5 Conclusion

I have argued that Weyl symmetry plays a substantial role in the derivation of the EFE in string theory. More precisely, the EFE follows from the hypothesis that the Weyl anomaly vanishes, and this hypothesis isn't empty. An adequate account of the Weyl anomaly requires a conception of symmetry that goes beyond the preserved–broken dichotomy found in Huggett and Vistarini's analysis and more broadly. I have indicated a replacement. On the alternative framing I have provided, the derivation of the EFE rests on the prohibition of gauge anomalies, and the justification of this prohibition should be further investigated. Leaving these details aside, I think Huggett and Vistarini's deflationary argument doesn't work. It finds Weyl invariance in every theory because it responds to failures of Weyl invariance by changing the theory under consideration. Some theories—indeed, most—are not Weyl invariant.

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