

On Miracles and Spacetime

James Read*

*Faculty of Philosophy, University of Oxford,
Woodstock Road, Oxford OX2 6GG, United Kingdom;*

*Department of Philosophy, University of Illinois at Chicago
1423 University Hall (M/C 267), 601 South Morgan Street
Chicago, IL 60607, United States*

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Abstract

What have recently been dubbed two ‘miracles’ of general relativity—(1) that all non-gravitational interactions are locally governed by Poincaré invariant dynamical laws; and (2) that, in the regime of experimental practice in which curvature effects may be ignored, the local Poincaré symmetries of the dynamical laws governing matter fields coincide with the local Poincaré symmetries of the dynamical metric field—remain unaccounted for in that theory. In this paper, I demonstrate that these two ‘miracles’ admit of a natural explanation in one particular successor theory to general relativity—namely, perturbative string theory. I argue that this point has important implications when considering both the ‘chronogeometricity’ (that is, the object in question being surveyed by rods and clocks built from matter fields) and spatiotemporal status of the dynamical metric field in both general relativity and perturbative string theory.

*james.read@philosophy.ox.ac.uk

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1 Introduction

But still we listen more gladly to others who hold out to us the hope that the more intimately we come to know the recesses of nature, or the better we are able to compare it with further aspects as yet unknown to us, the more simple shall we find its principles, and the further our experience advances the more harmonious shall we find it in the apparent heterogeneity of its empirical laws.

— Immanuel Kant, *Critique of Judgement*, p. 188.

Recent work in the philosophy of general relativity (GR) has drawn attention to two ‘miracles’—i.e., features “surprising or puzzling, the explanation of which remains outstanding” [45, p. 4, fn. 3]—in the foundations of the theory.¹ These ‘miracles’, labelled respectively **MR1** and **MR2**, can be stated as follows: [45, pp. 17-18]

MR1: All non-gravitational interactions are locally governed by Poincaré invariant dynamical laws.

MR2: The Poincaré symmetries of the dynamical laws governing non-gravitational fields in the neighbourhood of any point in the manifold coincide—in the regime in which terms featuring the Riemann tensor or its contractions can be ignored—with the symmetries of the dynamical metric field in that neighbourhood.

In this paper, after first recalling the details of **MR1** and **MR2**, I demonstrate that while these two results are indeed ‘miraculous’ in GR, they cease to be so in one particular successor theory of quantum gravity to GR, namely perturbative string theory. Not only is this a noteworthy observation in itself, but (I contend) it has important consequences when assessing both the ‘chronogeometricity’ (that is, the object in question being surveyed by rods and clocks built from matter fields²) and spatiotemporal status of the dynamical metric field in both general relativity and perturbative string theory.

To see the sense in which this is so, note two points. First, **MR1** and **MR2** constitute important conditions for the dynamical metric field in GR to acquire what is referred to in [7, 11, 45] as its ‘chronogeometric significance’—that is, for it to be surveyed by rods and clocks built from matter fields (see §3). In brief, the reason for this is that **MR2** ensures a surveying of the metric field by rods and clocks; **MR1**

¹The reader is urged not to be distracted by the nomenclature of ‘miracles’; cashing out exactly what qualifies as a ‘miracle’ is the central concern neither of [45] nor of the present paper. Rather, the above-quoted characterisation is fully sufficient for our purposes.

²Roughly, by ‘surveyed by rods and clocks built from matter fields’, I mean that rods and clocks built from matter fields can be used to read off intervals of the metric field; this notion is drawn from the work of Brown (see e.g. [7, ch. 9]), and is discussed in more detail in §3 below.

ensures the irrelevance of the constituents of rods and clocks.³ Second, if one follows Knox in defining ‘spacetime’ *functionally*, so that “the spacetime role is played by whatever defines a structure of local inertial frames” [28, p. 9], then, as we shall see, the obtaining of **MR1** and **MR2** is also an important condition for the metric field in GR to qualify as *spatiotemporal*. The fact that **MR1** and **MR2** are *contingent* in GR means that both chronogeometricity and spatiotemporality in GR is also contingent. Not so in perturbative string theory, where these two ‘miracles’ follow ineluctably from the formalism of the theory.

Consequently, both the chronogeometricity and spatiotemporality of the (non-fundamental) dynamical metric field in perturbative string theory have a stronger status than in GR. This prompts me to distinguish (A) the objects which play the functional role of spacetime *à la* Knox being fundamental in the theory, from (B) the objects which play the functional role of spacetime doing so necessarily rather than contingently. While GR exemplifies (A) but not (B), the situation in perturbative string theory is the reverse: the theory exemplifies (B), but not (A).

A final word of clarification. **MR1** and **MR2** were first articulated in [45], in which the so-called ‘dynamical approach’ to spacetime theories—first developed by Brown and Pooley [7, 9, 10]—was defended. Roughly speaking, this ‘dynamical approach’ involves two claims: (a) fixed fields, such as the Minkowski metric field η_{ab} of special relativity, are to be ontologically reduced to the symmetries of the dynamical equations governing matter fields (the dynamical view is, therefore, a modern form of relationalism—cf. [40, §6.3.2]); (b) ontologically autonomous metric fields, such as g_{ab} in GR, do not have their chronogeometric significance essentially (i.e., in all solutions of any theory in which they appear).⁴ Focussing on (b), advocates of one version of the opposing ‘geometrical approach’ to spacetime theories would state that ontologically autonomous metric fields, such as g_{ab} in GR, *do* have their ‘chronogeometric significance’ essentially. On such a geometrical approach, **MR1**

³Note, though, that **MR1** and **MR2** do not in themselves constitute sufficient conditions for such chronogeometricity, for one must also presuppose e.g. the existence of stable rods and clocks built from matter fields. For further discussion, see [45, pp. 15-16]. In that work, it is claimed that **MR1** and **MR2** instead constitute *necessary* conditions for chronogeometricity; however, it is perhaps more conservative to state that, with these above-mentioned auxiliary assumptions, they constitute *jointly* sufficient conditions for chronogeometricity.

⁴This latter point is what Butterfield calls in [12] ‘Brown’s moral.’

and **MR2** would not be ‘miraculous’ after all—for they would not be additional input assumptions in GR, but would, rather, follow ineluctably, simply in virtue of the presence of a metric field g_{ab} in this theory. However, in [43, 45], it was argued that this particular version of the geometrical approach is not viable—precisely because there exist problem cases for such a view, in which one has a metric field g_{ab} in one’s theory, but does not have the satisfaction of **MR1** and **MR2**. While in this paper I *do* assume that this particular version of the geometrical approach is untenable, I am otherwise agnostic on the dynamical/geometrical debate. Indeed, in [43], it was argued that there exist other, perfectly viable versions of the geometrical approach—but that **MR1** and **MR2** remain miraculous on such approaches (roughly speaking, these versions of the geometrical approach accept (b), but continue to reject (a)). Everything that I will say in this paper is compatible both with the dynamical approach, and with these more modest versions of the geometrical approach.⁵

The structure of this paper is as follows. In §2, I discuss **MR1** and **MR2**, and why these ‘miracles’ are merely contingent in GR. In §3, I discuss the chronogeometric and (for Knox) spatiotemporal status of the metric field in GR. In §4, I recall the relevant details of perturbative string theory, and demonstrate that, in this theory, **MR1** and **MR2** are necessary rather than contingent, and admit of a natural explanation in the string-theoretic context—thereby rendering them miraculous no more. In §5, I say a few more words on the logic underlying these ‘miracles’. In §6, I draw upon these results in order to compare the spatiotemporal status of the dynamical metric field in perturbative string theory, versus in GR. In §7, I discuss how Huggett’s arguments in [22] that ‘target space’ in perturbative string theory is not “phenomenal space” connect with these matters.

⁵While I do also discuss Knox’s inertial frame spacetime functionalism in this paper, it should be stressed that: (i) I do not commit myself, *à la* Knox, to inertial structure being the *sine qua non* of spatiotemporality; and (ii) commitment to this brand of spacetime functionalism (which certainly developed out of the dynamical approach—cf. [28]), is, in my view, best regarded as being *orthogonal* to a commitment in the dynamical/geometrical debate. For more on this latter point, see [35].

2 Two Miracles

In this section, I recall the content of **MR1** (§2.1) and **MR2** (§2.2) and their contingent status in GR, as a prelude to demonstrating in §4 how this situation changes on moving to perturbative string theory.

2.1 The First Miracle: Universal Local Poincaré Invariance

In order to understand **MR1**, one must recall some details regarding the dynamical equations of GR. The kinematically possible models (KPMs) of GR are picked out by triples $\langle M, g_{ab}, \Phi \rangle$,⁶ where M is the a four-dimensional differentiable manifold; g_{ab} is a generic Lorenzian metric field on M ; and Φ is a placeholder for matter fields in the theory.⁷ Dynamically possible models (DPMs) of GR are those KPMs the geometrical objects of which obey certain specified dynamical equations: the Einstein field equations,

$$G_{ab} := R_{ab} - \frac{1}{2}g_{ab}R = 8\pi T_{ab}, \quad (2.1)$$

i.e. the dynamical equation associated with g_{ab} , *plus* dynamical equations for the matter fields.⁸ (Here, T_{ab} is the stress-energy tensor associated with the matter fields Φ .)

⁶Throughout this paper, abstract (i.e. coordinate-independent) indices are written in Latin script; indices in a coordinate basis are written in Greek script; and (where appropriate) the Einstein summation convention is used. I set $c = G_N = 1$.

⁷In this paper, I mean by ‘matter fields’, or ‘non-gravitational fields’, those for which there exists an associated stress-energy tensor, and by ‘gravitational fields’ those for which there exists no such stress-energy tensor—this distinction is in the spirit of [31]. In the context of general relativity, this means that the metric field is identified as a gravitational field, whereas all other fields typically of interest (e.g. Klein-Gordon fields, electromagnetic fields, etc.) are matter fields. There exist here issues regarding the possibility of defining a stress-energy tensor associated with the metric fields in general relativity—see [14] for a proof against this possibility, and e.g. [20, 29, 44] for related discussion.

⁸Strictly, independence of these dynamical equations from the Einstein field equations depends on the case in question—for discussion of this point, see [7, §9.3] and [36, §20.6].

Ab initio, there is no restriction on the possible dynamical equations for matter fields in GR. However, one often restricts consideration to equations which are *minimally coupled*. The prescription for constructing such minimally coupled equations is the following:

Minimally coupled dynamical equations for matter fields in GR are constructed from dynamical equations for matter fields featuring coupling to a fixed Minkowski metric field η_{ab} and no curvature terms, by replacing all instances of η_{ab} with a generic Lorentzian metric field g_{ab} , and replacing all instances of the torsion-free derivative operator compatible with η_{ab} with the torsion-free derivative operator compatible with g_{ab} .

Sometimes, this is referred to as the ‘comma-to-semicolon’ prescription. To make further progress in understanding minimal coupling, recall now some basic facts regarding pseudo-Riemannian geometry. In a coordinate basis $\{e_\mu\}$, the *connection components* $\Gamma^\mu_{\nu\rho}$ associated to a derivative operator ∇_a are defined by $\nabla_\rho e_\nu =: \Gamma^\mu_{\nu\rho} e_\mu$. Then, at any $p \in M$, we can choose *normal coordinates*, such that $\Gamma^\mu_{(\nu\rho)}(p) = 0$ in those coordinates; for a torsion-free derivative operator, we can in fact choose normal coordinates such that $\Gamma^\mu_{\nu\rho}(p) = 0$. (Note that the connection components away from p will in general *not* vanish.) If the unique torsion-free, metric compatible derivative operator is used, then in normal coordinates we also have $\partial_\rho g_{\mu\nu}(p) = 0$, and we may further restrict to the subclass of normal coordinates at p such that $g_{\mu\nu}(p) = \text{diag}(-1, 1, 1, 1)$.

Minimally coupled dynamical equations for non-gravitational fields in GR take a Poincaré invariant form in normal coordinates at any $p \in M$ —see [45, §A]. Indeed, the work of that paper demonstrates that *any* dynamical equation featuring coupling to (a) g_{ab} ; (b) its associated curvature tensor; (c) the matter fields in the theory, will also take Poincaré invariant form in normal coordinates at any $p \in M$ —even if such a dynamical equation is not obtained via the minimal coupling prescription. On the other hand, many dynamical equations are possible in GR which are *not* locally Poincaré invariant—for example, there is nothing to rule out dynamical equations for non-gravitational fields in GR which are locally *Galilean* invariant.⁹

⁹Presumably, KPMs here would be picked out by tuples $\langle M, g_{ab}, t_{ab}, h^{ab}, \nabla_a, \Phi \rangle$, where g_{ab} is a

The question posed by **MR1**, then, is: why should the dynamical equations for non-gravitational fields be such that all are Poincaré invariant? Given that other coupling schemes are possible in which such local Poincaré invariance does not obtain, that such *is* the case constitutes a ‘free lunch’. Within GR, it does not appear that this result can be delivered without further input assumptions—e.g., the restriction to minimally coupled dynamical equations. But in any case, satisfaction of **MR1** must ultimately simply be *postulated*; it does not appear that any deeper explanation of this ‘miracle’ is possible, within the framework of GR.

2.2 The Second Miracle: Metric Symmetry Coincidence

Turn now to an explication of **MR2**. First, recall that, in normal coordinates at p , $\partial_\rho g_{\mu\nu}(p) = 0$, and one may further restrict to the subclass of normal coordinates such that $g_{\mu\nu}(p) = \text{diag}(-1, 1, 1, 1)$; however, higher derivatives of the metric field do not necessarily vanish at p . Second, assume that we are dealing with matter fields which *do* all satisfy Poincaré invariant dynamical laws at any $p \in M$. Third, note that such dynamical equations for matter fields may still feature curvature terms, even at p (for an extended discussion of this point, see [45, §§2-3]). Since higher derivatives of the metric field do not vanish at p , it would be misleading to state that the metric field g_{ab} ‘reduces’ to the Minkowski metric field η_{ab} at p . However, since non-vanishing higher derivatives of the metric field are associated with non-vanishing derivatives of connection components, which in turn are associated with non-vanishing curvature, we can say that, *in the regime in which terms featuring derivatives of connection components can be neglected* (whether this is so depends upon the strength of one’s experimental apparatus relative to the strength of the ‘tidal gravitational forces’ represented by these terms—for detailed discussion, see [45, §3.2]), the metric field g_{ab} *does* recover the form of a Minkowski metric η_{ab} at p ,

generic Lorentzian metric field satisfying (2.1); $\langle M, t_{ab}, h^{ab}, \nabla_a \rangle$ is the structure of ‘classical space-time’ (in the sense of [33, ch. 4]) necessary to write down Galilean invariant laws for non-gravitational fields; Φ are the matter fields, which obey Galilean invariant dynamical laws; and the stress-energy tensor T_{ab} in (2.1) is that associated with the Φ . In this scenario, g_{ab} dynamically evolves in response to the stress-energy content of the Φ , but the Φ are not so sensitive to g_{ab} —for this metric field does not feature in their associated dynamical equations. For this reason, we seem to have here instantiation of a strange, *reverse* violation of the action-reaction principle, since the metric field is responsive to matter fields but not vice versa (for further discussion of the action-reaction principle, see [7, 8]).

and the dynamical laws governing matter fields reduce to a Poincaré invariant form with no curvature terms, at p .

What this means is that, in such a regime, the symmetries of the metric field at $p \in M$ —i.e., the coordinate transformations under which the metric at p retains the same form—coincide with the symmetries of the dynamical laws governing matter fields at p —i.e., the coordinate transformations under which these dynamical equations retain the same form. That such *is* the case is a statement of **MR2**—again, this constitutes a contingent ‘miracle’ in GR, for it is in principle possible that all matter fields are governed by dynamical equations with the same local symmetries, which nevertheless do not coincide with those local symmetries of the metric field—the above example of all matter fields obeying Galilean invariant local dynamical laws constitutes just such a case.

3 Chronogeometry and Spacetime

Why is it that intervals of the metric field g_{ab} in GR can be read off from rods and clocks constructed from matter fields? An important condition for the metric field of GR to have ‘chronogeometric significance’ in this sense is that **MR1** and **MR2** hold. To see this, suppose that these principles hold, so that all dynamical laws for matter fields have the same symmetries, which locally coincide (in the appropriate regime) with those of the metric field. Now consider a stable rod or clock built from matter fields, in a particular frame.¹⁰ If such a rod or clock is a ‘good’ rod/clock in that frame (if, one might say, it is ‘ideal’ in that frame), then it can be used to measure intervals as given by the metric field, in that frame. But now, since **MR1** and **MR2** are assumed to hold, in a Poincaré-transformed frame such is *still* the case, for metric and dynamical equations for matter fields transform in the same way—meaning that this stable rod/clock yields a robust means of measuring such metric intervals in all frames related by Poincaré transformations.¹¹

¹⁰Recall from footnote 3 that the existence of such stable rods and clocks constitutes an important additional condition for the metric field of GR to have chronogeometric significance.

¹¹By contrast, suppose that one has a rod/clock built from matter fields which obey Galilean invariant laws; then, in a Poincaré-transformed frame, that rod/clock would in general not obey dynamical

Since **MR1** and **MR2** are contingent in GR, so too is this chronogeometric significance of the metric field: if different matter fields obeyed different dynamical laws, then the metric field would not have (universal) local chronogeometric significance; if the symmetries of dynamical laws for matter fields did not coincide locally with those of the metric field, then, again, the metric field would not have local chronogeometric significance. Indeed, not only is the chronogeometric significance of the metric field in GR a contingent matter, but so too is its *spatiotemporal status*, on a spacetime functionalist approach such as that of Knox [26–28]. Recall again Knox’s slogan: “I propose that the spacetime role is played by whatever defines a structure of local inertial frames” [28, p. 9]. The idea is to functionally *define* spacetime as any structure which itself picks out a structure of local inertial frames. In turn, Knox gives the following characterisation of inertial frames: [26, p. 348]

In Newtonian theories, and in special relativity, inertial frames have at least the following three features:

1. Inertial frames are frames with respect to which force free bodies move with constant velocities.
2. The laws of physics take the same form (a particularly simple one) in all inertial frames.
3. All bodies and physical laws pick out the same equivalence class of inertial frames (universality).

Any structure which picks out a “structure of local inertial frames”, i.e. a structure of local frames which satisfy these properties (initially identified as significant in the Newtonian/special relativistic context) qualifies definitionally as ‘spacetime’, for Knox. But note that, in GR, **MR1** and **MR2** precisely guarantee that this field be considered spatiotemporal, in this sense. The reason is that, locally, the symmetries of the dynamical metric field coincide with those of the dynamical equations governing matter fields; in any frame in which these dynamical equations take their simplest form, the metric field itself takes the form $\text{diag}(-1, 1, 1, 1)$. Thus, the metric field picks out a structure of local inertial frames.

equations of the same form—and so would in general not be a ‘good’ rod/clock in that frame.

4 Perturbative String Theory

The above in hand, I turn now to perturbative string theory—one candidate quantum theory of gravity.¹² Classically, a string can be regarded as a special case of a p -brane, which is an object with p dimensions and tension $T_p = 1/(2\pi l_S^2)$, where l_S is the so-called ‘string length’. The classical motion of a p -brane extremises the $(p+1)$ -dimensional volume V that it sweeps out in a D -dimensional ‘target space’.¹³ Thus, there is a p -brane action given by $S_p = -T_p V$. In the case of the fundamental string, which has $p = 1$, V is the area of the string worldsheet and the action is called the *Nambu-Goto action* [5, p. 10]

$$S_{\text{NG}} = -\frac{1}{2\pi l_S^2} \int_{\Sigma} d^2\sigma \left[-\det_{ab} \left(\frac{\partial X^a}{\partial \sigma^a} \frac{\partial X^b}{\partial \sigma^b} \eta_{ab} \right) \right]^{1/2}, \quad (4.1)$$

where Σ denotes the string worldsheet; the functions $X^a(\sigma, \tau)$ describe the target space embedding of the worldsheet; $\tau \equiv \sigma^0$ and $\sigma \equiv \sigma^1$ are coordinates on the worldsheet (the parameter τ is the worldsheet ‘time’ coordinate, while σ parametrises the string at a given worldsheet ‘time’); $d^2\sigma = d\tau d\sigma$; and Fraktur script denotes worldsheet indices. Classically, the Nambu-Goto action is equivalent to the string sigma-model action, also known as the *Polyakov action* [5, p. 13]

$$S_{\text{P}(\eta)} = -\frac{1}{4\pi l_S^2} \int_{\Sigma} d^2\sigma \sqrt{-h} h^{ab} \eta_{ab} \partial_a X^a \partial_b X^b, \quad (4.2)$$

where $h_{ab}(\tau, \sigma)$ is an auxiliary¹⁴ worldsheet metric with inverse $h^{ab}(\tau, \sigma)$; and $h := \det h_{ab}$. The Euler-Lagrange equation for h_{ab} can be used to eliminate it from (4.2)

¹²String theory goes further than other theories of quantum gravity (e.g. loop quantum gravity), since it purports to realise the more ambitious goal of unifying all four fundamental forces into one (quantum mechanical) framework. That is, string theory purports to be a ‘theory of everything’, whereas other quantum gravity theories need not do so.

¹³Recently, in light of so-called *dualities* in string theory (see e.g. [47]), Huggett has argued that “phenomenal spacetime” is *not* equivalent to target space [22]—see §7 for an extended discussion of this point.

¹⁴In the sense that h_{ab} is a *new* variable, *a priori* independent of the pullback of the target space metric to the worldsheet.

and recover (4.1).¹⁵ Quantum mechanically, instead of eliminating h_{ab} via its classical field equations, one should perform a path integral, using standard machinery to deal with the local symmetries and gauge fixing.¹⁶ Doing this, one finds that there is a conformal anomaly¹⁷ unless the target space dimension is $D = 26$. Analogous analysis for superstrings (i.e. strings for which supersymmetry is added—either on the worldsheet as in the so-called *RNS sector*, or to the background target space as in the *GS sector*)—gives the critical dimension $D = 10$ [2, p. 7]. (In this section, I focus for simplicity solely upon bosonic string theory.)

For a closed string, one imposes periodicity in σ . Choosing its range to be π , one identifies both ends of the string $X^a(\tau, \sigma) = X^a(\tau, \sigma + \pi)$. After quantising and defining suitable ladder operators (see e.g. [2, p. 53]), one can act on the ground state of the string with raising operators to study its spectrum. For the closed string, there exist three distinct first excited states, denoted f_{ab} , B_{ab} , and Φ . f_{ab} —suggestively christened the *graviton*—is symmetric and traceless in its target space indices, and transforms under $SO(D - 2)$ as a massless,¹⁸ spin-two particle. B_{ab} transforms under $SO(D - 2)$ as an antisymmetric, second-rank tensor. The trace term Φ is a massless scalar, which is called the *dilaton* [2, p. 53].

In (4.2), I considered only a fixed, flat background η_{ab} . One can analyse more general possibilities by introducing the fields f_{ab} , B_{ab} , and Φ into the worldsheet action as background fields. The appropriate action is then [5, p. 429]

¹⁵Varying (4.2) with respect to h^{ab} , we obtain the equation of motion $\eta_{ab}\partial_a X^a\partial_b X^b - \frac{1}{2}h_{ab}h^{cd}\eta_{cd}\partial_c X^a\partial_d X^a = 0$; back-substitution then returns (4.1).

¹⁶I.e. path integral quantisation *à la* Faddeev-Popov—see e.g. [5, §3.4].

¹⁷An *anomaly* arises when a symmetry of a classical theory is not manifested in the associated quantum theory. In more technical language, though the classical action is invariant under the symmetry, the associated path integral measure—used to define the quantum theory—is not. The *conformal anomaly* arises on quantisation of classical string theory, and breaks the conformal invariance of the string worldsheet. This anomaly manifests itself as an extra term in the *Virasoro algebra*, which comprises the generators of the conformal group of the string worldsheet; this extra term vanishes—thereby circumventing the anomaly—only in the case of target space dimension $D = 26$ for the bosonic string, or $D = 10$ for the fermionic string. For a philosophically-oriented introduction to anomalies, see [21, §4].

¹⁸For how the masses of such states are determined, see e.g. [2, p. 53].

$$S_{P(\eta+f,B,\Phi)} = -\frac{1}{4\pi l_S^2} \int_{\Sigma} d^2\sigma \sqrt{-h} \left(h^{ab} \partial_a X^a \partial_b X^b (\eta_{ab} + f_{ab}(X)) \right. \\ \left. + \epsilon^{ab} \partial_a X^a \partial_b X^b B_{ab}(X) + l_S^2 \Phi R(h) \right), \quad (4.3)$$

where conventions are such that $\epsilon^{ab} = \pm 1/\sqrt{-h}$. Thus, in perturbative string theory there is a move from finding the excited states of strings, to treating the associated quantum fields as *background fields* in which the dynamics of further strings can be analysed.¹⁹

Consider a string propagating in background fields, as per (4.3). This action may be regarded as that of a conformal field theory (of the X^a fields) on the string world-sheet. It turns out, remarkably, that this conformal invariance *requires for consistency* that the background fields be dynamically coupled together in the Einstein field equations, plus higher order corrections.²⁰ This follows since the beta function²¹ associated to each of the background fields (each such background field being treated as a coupling parameter in the worldsheet conformal field theory) can be computed perturbatively; conformal invariance then requires these beta functions to vanish. For example, for the closed bosonic string with background fields g_{ab} , B_{ab} and Φ , and $D = 26$, one finds the following one-loop beta functions associated to each of the background fields, having performed a conformal rescaling $g_{ab} \rightarrow e^{\Phi/6} g_{ab}$,²² [13, p. 597]

¹⁹Attempting to make sense of this move is a philosophically nuanced task—see e.g. [21, 23, 42]. In this paper, I pass over the subtleties which arise in this vicinity, and simply *assume* that it is legitimate to treat excited states of strings as background fields, as in (4.3).

²⁰For original sources on this result, see e.g. [13]; for textbook discussion, see [19, pp. 167ff.]; for a presentation in the philosophical literature, see [23, §3].

²¹In quantum field theory, a *beta function* $\beta(g) = \partial g / \partial \log(\mu)$ encodes the dependence of a coupling parameter g upon an energy scale μ . Since a conformal field theory does not depend upon an energy scale, the beta function associated to its coupling parameters must vanish. See e.g. [38, 51] for further details.

²²Here, H_{abc} is the generalised Faraday tensor associated to B_{ab} —so that, in the notation of differential forms, $H = dB$.

$$G_{ab} = 8\pi T_{ab}, \quad (4.4)$$

$$\nabla_a \nabla^a \Phi + \frac{1}{6} e^{-\Phi/3} H^2 = 0, \quad (4.5)$$

$$e^{-\Phi/3} \nabla_a H^a_{bc} = 0, \quad (4.6)$$

with

$$T_{ab} = \frac{1}{4} \left(H_{acd} H^{cd}_b - \frac{1}{6} g_{ab} H_{abc} H^{abc} \right) e^{-\Phi/3} + \frac{1}{6} \left(\nabla_a \Phi \nabla_b \Phi - \frac{1}{2} g_{ab} \nabla_a \Phi \nabla^a \Phi \right), \quad (4.7)$$

and G_{ab} the (26-dimensional analogue of the) Einstein tensor. Analogous results for the *heterotic superstring* (in which there also exists a one-form background field A_a) yield more familiar coupled Maxwell-Einstein systems (see [13, §4]).²³

Allowed background fields in perturbative string theory are restricted to those which correspond to the excited states of strings; one such field ($g_{ab} := \eta_{ab} + f_{ab}$, where f_{ab} is the graviton field) corresponds to a Lorentzian metric field; the other fields are identified as non-gravitational fields.²⁴ Conformal invariance of the string worldsheet metric—itsself required for the consistency of the theory, and hence not an input assumption therein—guarantees that g_{ab} satisfies the Einstein field equations, and that all matter fields obey Poincaré invariant dynamical laws. Thus, in perturbative string theory, *unlike* GR, one is *not* free to choose matter field couplings; rather, the dynamical equations governing the behaviour of matter fields on target space are an ineluctable consequence of the formalism. In other words, the space of DPMs for the background fields in a given string theory has a *fixed* set of dynamical equations for matter fields. For this reason, **MR1** is, in string theory, a necessary principle. Indeed, such is also the case for **MR2**: it is again a consequence

²³These results are not presented explicitly in order to avoid technical subtleties regarding superstring theories; the above will suffice for our purposes.

²⁴See footnote 7, where a distinction was drawn between gravitational and non-gravitational fields.

of the formalism that (in the relevant regime) the local Poincaré symmetries of the metric field g_{ab} coincide with those of the dynamical equations governing matter fields. Thus, on moving to a successor theory to GR, **MR1** and **MR2** both *cease to be miraculous*.

5 Logic of the Miracles

It is worth taking some time to reflect upon the logic of the status of **MR1** and **MR2** as they appear in GR, versus as they appear in perturbative string theory. In the former case, **MR1** must simply be postulated as an input assumption; **MR2** then follows once it is noted that there also exists in GR an autonomous Lorentzian metric field, the symmetries of which happen to coincide with the symmetries of the dynamical laws governing matter fields in the theory, assuming **MR1**. Thus, the situation in GR is as follows:

$$\begin{array}{ccc} & & \mathbf{MR1} \\ \mathbf{MR1} + \text{Lorentzian } g_{ab} \text{ in KPMs} & \Rightarrow & \mathbf{MR2} \end{array}$$

Contrast this with the case of perturbative string theory. In that case, one posits a Minkowski metric field η_{ab} on the target space, but having done so **MR1** and **MR2** follow as consequences of the formalism, with no additional input assumptions. Thus, the situation in perturbative string theory is as follows:

$$\eta_{ab} \text{ in KPMs} \Rightarrow \mathbf{MR1} + \mathbf{MR2}$$

Clearly, this is not to say that there are *no* input assumptions in perturbative string theory at all; it is, rather, to say two things: (a) the *number* of input assumptions in perturbative string theory versus GR is reduced; (b) what were two unex-

plained coincidences in GR (*viz.*, **MR1** and **MR2**) admit of a deeper, overarching explanation in the case of perturbative string theory.²⁵

To put matters differently, in the case of GR, **MR2** consists of two points: (i) a restatement of **MR1**, and (ii) the observation that one is dealing with a Lorentzian metric field g_{ab} in the KPMs of the theory. Clearly, it is (ii) that is the extra component of **MR2**, over and above **MR1**. But note that this extra assumption is analogous to the sole input assumption in string theory, that one is dealing with a Minkowski metric field η_{ab} in the KPMs of the theory.²⁶ In the case of string theory, this extra assumption suffices, ultimately, to account for **MR1**, whereas such is *not* the case in GR. Thus, while in GR there exist two input assumptions, in string theory there exists only one.

A final word on the so-called ‘landscape problem’ of string theory.²⁷ Recall that the target space of superstring theory has $D = 10$; in order to recover the phenomena of our four-dimensional world, six of these dimensions are typically taken to be ‘compactified’—loosely, to be associated with periodic dimensions of small radius (smaller than any foreseeable radius of empirical detectability)—see e.g. [2, ch. 9]. In order to preserve supersymmetry, these six dimensions are standardly taken to be those of a Calabi-Yau threefold.²⁸ However, many inequivalent choices for a Calabi-Yau threefold are possible; each different choice leads ultimately—via the analysis presented in §4—to different background fields on the target space, obeying different dynamical equations. Given this, one might ask in response to the

²⁵Setting up the logic of the two miracles in this way raises an interesting question: would it be the case that one would recover the analogues of **MR1** and **MR2** in string theories with the structure of (e.g.) *Galilean* spacetime on the target space (cf. e.g. [18])? The study of alternative target space metric fields may raise interesting complications in this vicinity—for example, need it be the case that conformal invariance of the string worldsheet is necessary for anomaly cancellation, when one has such alternative target space structure? If not, then it may not be so easy to argue—for these alternative target spaces—that the dynamics of background fields is fixed, and therefore that (the relevant equivalents of) **MR1** and **MR2** are consequences of the formalism of the theory. My thanks to an anonymous referee for helpful discussions on this matter.

²⁶In fact, this latter assumption is somewhat stronger, for in the case of string theory one assumes a specific form of the target space metric field (namely, that it is the Minkowski metric), whereas in GR one assumes only that this metric field be Lorentzian. Many thanks to Nick Huggett for useful discussion on this point.

²⁷For a foundational paper on this problem in string theory, see [50].

²⁸Calabi-Yau manifolds are compact Kähler manifolds with vanishing first Chern class—see e.g. [2, §9.2] for further background. ‘Threefold’ here means three complex dimensions—clearly, this corresponds to six real dimensions.

present paper: “How can string theory possibly be understood as fixing the dynamical equations of matter fields?” My response to this is the following. I do not deny that, in perturbative string theory, there are many inequivalent choices of ‘vacua’—that is, background target space manifolds, to which one applies the machinery of perturbative string theory as in §4. Moreover, I do not deny that, ultimately, it would be preferable to have some *non-perturbative* version of string theory, which does not involve fixing a background manifold for target space.²⁹ Rather, what I claim is the following: *given* a certain choice of target space manifold—namely, a Lorentzian manifold with Minkowski metric field η_{ab} —the allowed dynamical equations for matter fields are fixed; thus, one is not in general free to choose these equations independently of one’s assumptions about the metric field on target space, as one is in GR. For the particular choice of a Minkowski metric field η_{ab} on target space, then, I have demonstrated in this paper that **MR1** and **MR2** are fixed—so, unlike GR, they cease to be input assumptions in the theory.³⁰

6 Emergent Spacetime

The observation that **MR1** and **MR2** *must* hold in perturbative string theory has important implications for the emergence of spacetime in quantum gravity. Drawing upon Knox’s discussion at [26, §1], distinguish two senses in which the ‘emergence of spacetime’ might be understood. On the former, such emergence arises whenever an object (or collection of objects) in a given theory is not fundamental—in the sense of its not being defined at the level of KPMs of the theory, but rather being constituted out of those fundamental objects specified in the KPMs, in a possibly limited range of physical circumstances—but which nevertheless plays the functional role of spacetime as characterised in §3. Knox claims (correctly) that one need not look to theories of quantum gravity to find illustrations of this phenomenon—for example, teleparallel gravity (an alternative theory of classical gravity to GR, in

²⁹For more on this, see §6 below.

³⁰Thus, *even if*, for particular choices of the target space manifold, one witnesses violations of **MR1** and **MR2**, this would not invalidate my point that, *given* a choice of target space manifold, the dynamics of background fields are fixed—so that satisfaction (or violation) of **MR1** and **MR2** may not be freely chosen. My thanks to an anonymous referee for discussion on this matter.

which curvature degrees of freedom are put into torsion in the connection—for Knox’s discussion of this theory, see [25,26]; for further philosophical details, see [42, ch. 4]) seems to instantiate just such a point, for in this theory those spatiotemporal degrees of freedom are composite objects, built from fundamental fields specified at the level of KPMs.³¹

The second sense in which spacetime can be emergent is the following: a given object’s playing the functional role of spacetime *à la* Knox is not a necessary consequence of the formalism, in the sense that this object (fundamental or not) plays that functional role in all models of the theory, but is rather contingent, in the sense that such an object plays the spacetime role in some models of the theory, but not others. Again, Knox correctly observes at [26, §1] that one does not need to look to theories of quantum gravity for illustrations of this case—GR, for the reasons outlined in §3, constitutes a clear example of this latter sense of emergence.

Though GR exemplifies the emergence of spacetime in the latter sense, it does *not* exemplify it in the former sense, for in this theory it is the *fundamental* metric field g_{ab} which plays the functional role of spacetime. The situation is very different in perturbative string theory, in which the object playing the functional role of spacetime is the composite g_{ab} , built out of the Minkowski metric field η_{ab} and graviton field f_{ab} , the latter itself a coherent state of strings in an excited graviton mode.³² Since this object is non-fundamental, it qualifies as emergent spacetime in the former sense; however, it does *not* qualify as emergent spacetime in the latter sense, for this object *necessarily* plays the functional role of spacetime, in all models of the theory.³³ In other words, GR and perturbative string theory are, in a certain sense, ‘dual’ *vis-à-vis* emergent spacetime in these theories, for while spacetime in GR is *fundamental but contingent*, spacetime in perturbative string theory is *non-fundamental but necessary*.

³¹Another classical theory which exemplifies the emergence of spacetime in this sense is Bekenstein’s bimetric TeVeS (‘Tensor-Vector-Scalar’) theory, presented in [3, 4]. As discussed in [7, §9.5.2], in this theory the metric field which is surveyed by rods and clocks, the conformal structure of which is traced by light rays, and the geodesics of which correspond to the motion of free bodies, is not the ‘fundamental’ metric field g_{ab} , but rather a less ‘fundamental’ metric field $\tilde{g}_{ab}$, constructed from the other matter fields in the theory [7, p. 174].

³²For some discussion of the extent to which the η_{ab} field can be eliminated in perturbative string theory, see [42, ch. 8].

³³Here, I take ‘the theory’ to refer to perturbative string theory, with η_{ab} on the target space.

Now, it is worth noting that perturbative string theory does not, of course, exhaust the string-theoretic picture. While the above story might be correct for perturbative string theory, it might *not* hold in general for a *non-perturbative* formulation of the theory. For example, it has been conjectured that some solutions of ‘M-theory’—perhaps the best-known candidate for a non-perturbative formulation of string theory—will not rely, fundamentally or otherwise, upon spatiotemporal concepts.³⁴ If that is correct, then it may be that the idea of spacetime being necessary in non-perturbative string theory does not hold water, unlike the situation in the perturbative case. Though these observations do not undermine the claims of the current paper, they certainly do deserve to be flagged.

7 Dualities

Up to this point, I have taken it for granted that, on a spacetime functionalist approach, the target space metric field $g_{ab} := \eta_{ab} + f_{ab}$ represents spacetime. However, at this juncture I must acknowledge recent work by Huggett, in which it is argued that this target space metric field cannot represent “phenomenal spacetime”, i.e., “the geometrical space we take ourselves to experience in the everyday, including the experience of three large dimensions” [22, p. 7]. The reason for this has to do with the phenomenon of *T-duality* in string theory.³⁵

T-duality is just one of a broad class of ‘dualities’ in theoretical physics in general, and string theory in particular. For the purposes of this paper, a *duality* can be taken to be a mapping between the spaces of DPMs of two theories, such that any two models related by that mapping are empirically equivalent (in turn, the empirical equivalence of two models can be defined in terms of their agreeing on all *empirical substructures*, in the sense of van Fraassen [54, pp. 67ff.]).³⁶ In the case of T-duality, models of perturbative string theory on a target space product manifold $M \times S^1$ with

³⁴The question of whether M-theory may or may not deploy spatiotemporal concepts was the sixth open question on string theory presented by Strominger at the *Strings 2014* meeting—cf. [52]. For some philosophical discussion of this matter, see [34, p. 14].

³⁵See [24, §2.4] for related discussion, and [34] for an extension of Huggett’s arguments to other string-theoretic dualities.

³⁶For more on the definition of dualities, see e.g. [16, 17, 41, 46, 47].

radius of the periodic dimension R are found to be dual to models of perturbative string theory on the target space product manifold $M \times S^1$ with radius of the periodic dimension equal to l_S^2/R (see e.g. [2, ch. 6]).

Within the literature on dualities, some (e.g. [47, 49]) have taken it that duality-related models should be regarded *ab initio* as representing the same possible world (i.e. as being not just empirically equivalent, but also *physically* equivalent); others (e.g. [16, 46]) have resisted this conclusion. Assuming that one endorses the former approach, then, since duality-related models often have very different mathematical structures, there exists a burden to account for the common ontology which such models are taken to represent; one popular option is to seek an explication in terms of the mathematical structure *common* to both models.³⁷ Not only does Huggett fall into the former interpretative camp—that is, he takes it that duality-related models should be regarded *ab initio* as representing the same physical state of affairs³⁸—but, in addition, he maintains that the physical state of affairs represented by such duality-related models should be explicated in terms of their common structure [22, §2.3].

In light of this, Huggett’s position on the interpretation of T-dual models of perturbative string theory can be summarised as follows. While “[i]t is natural when first introduced to string theory, to think that target space is simply the same space we ordinarily experience, or at least space as conceived in contemporary physics” [22, p. 8], ultimately “T-duality shows that a definite radius for target space ... [is] not physical, but only [an artefact] of the representation” [22, p. 10], for this radius is not part of the structure common to the two models under consideration. Since this aspect of the target space metric field is unphysical, for Huggett, it follows that the target space metric field g_{ab} cannot always represent phenomenal spacetime; rather, what *does* represent phenomenal spacetime must be constructed from the shared structure of the T-dual models. In fact, following arguments by Brandenberger and Vafa [6], Huggett maintains that phenomenal spacetime always coincides with the

³⁷For a detailed discussion of this option and others, see [30, 41].

³⁸At least for theories which are supposed to constitute descriptions of the entire universe—see [22, §2.2], and [16, §3.1] for further discussion.

target space of the *larger* radius between two T-dual models.^{39,40}

If Huggett is correct here, then does this pose a problem for our verdict that spacetime is ‘non-fundamental but necessary’ in perturbative string theory? On the one hand, one might be inclined to answer this question in the negative, for (one might say) there is nothing to stop an object in a theory playing the functional role of spacetime, while at the same time not being a candidate for representing phenomenal spacetime, or ultimately being regarded as ‘surplus structure’. Nevertheless, one might feel that there is more to be said here, and that an affirmative answer to the above question is also possible. In particular, Huggett’s discussion of phenomenal spacetime as being that construct which is surveyed by *macroscopic* rods and clocks might make one think that whatever plays the functional role of spacetime should *also* satisfy this criterion. In that case, one should be disturbed by the verdict of Knox’s spacetime functionalism, that what plays the functional role of spacetime need not coincide with phenomenal spacetime in this sense.

If one buys this reasoning, then one might read this discussion of T-duality as highlighting that Knox’s functional definition of spacetime as that structure which “defines a structure of local inertial frames” cannot be *sufficient* for the identification of spatiotemporal structure, for if what qualifies functionally as spacetime must *also* qualify as phenomenal spacetime, then this object should, as Huggett notes, also be surveyed by macroscopic rods and clocks. Only one of the T-dual target space metric fields does as much; thus, with this extra criterion imposed, our above conclusions must be revised, and we should—it appears—conclude that spacetime in perturbative string theory is neither fundamental nor necessary, in the senses of

³⁹Note, though, that Huggett does not maintain that phenomenal spacetime *just is* the target space of the larger radius, for the model under consideration with the *smaller* radius will also (of necessity, in light of the duality) manifest the mathematical structure needed to represent phenomenal spacetime, while not exhibiting fundamentally a target space of large radius. Thanks to Nick Huggett for stressing this point.

⁴⁰An anonymous referee has rightly pointed out to me that there is some tension between this claim that, in the case of T-dual models, phenomenal spacetime always coincides with the target space of the larger radius, and Huggett’s earlier claim that phenomenal spacetime is “the geometrical space we take ourselves to experience in the everyday, including the experience of three large dimensions” [22, p. 7]—for the former, at least as presented above, will have more than three “large” spacetime dimensions. Here is one way to resolve the tension: by “phenomenal spacetime”, Huggett means the classical structures which represent the spatiotemporal data which we *would* experience, *were we* to be ‘embedded’ in the model under consideration. If that is right, then (e.g.) having three dimensions is not, in fact, necessary for a given structure to qualify as “phenomenal spacetime”.

§6.

In fact, however, the situation on these matters is even more subtle, for one can argue that, in any model of T-dual perturbative string theories, one can (in light of the formal duality map) always identify *some* structure corresponding to phenomenal spacetime. If that is correct, then one may be able to argue that every model of T-dual string theory *does* possess some structure corresponding to spacetime, even on the above-proposed revised Knoxian account. In that case, spacetime in T-dual perturbative string theories would be non-fundamental but necessary after all. Indeed, it would be *doubly* non-fundamental, for not only are the target space metric fields g_{ab} composite objects, but, in addition, what represents (phenomenal) spacetime should not (as Huggett stresses) be identified straightforwardly with those objects.⁴¹

Let me recapitulate the dialectic of this section. First, I observed that Knox's original functionalist approach to spacetime would, it appears, identify the target space metric fields g_{ab} in models of perturbative string theory as being spatiotemporal. On that view, spacetime in perturbative string theory is non-fundamental, but necessary. I then observed (following Huggett), however, that target space metric fields need not be identified as spatiotemporal, for they need not be surveyed by macroscopic rods and clocks. If one augments Knox's functional definition of spacetime to include also this criterion, then (it appears—so long as one continues to focus one's attention upon the target space metric fields) that spacetime in perturbative string theory is non-fundamental and contingent. Finally, however, I reflected further upon phenomenal spacetime. Since every model of T-dual perturbative string theory possesses sufficient mathematical structure to represent phenomenal spacetime (in light of the existence of the duality map—so even those models of perturba-

⁴¹Let me mention one issue which might complicate the dialectic yet further. Consider T-dual models in which the radius is close to the 'minimum length', so that $R \approx l_s^2/R$. In this case, is it obvious that photons, *à la* Brandenberger and Vafa, should be understood as reading off intervals as given by *either* target space metric field? In this regard, cf. [34], in which it is argued that, in some models of perturbative string theory related by more complicated dualities such as 'mirror symmetry' (which is, essentially, a generalisation of T-duality to topologically inequivalent manifolds—see [48] for a philosophically-oriented introduction), there need not be *any* structure in the models corresponding to phenomenal spacetime (so that, arguably, such models are 'unphysical'). If this is correct, then spacetime in string theory would once again become non-fundamental and contingent. My thanks to an anonymous referee for helpful discussion on this issue.

tive string theory with *small* target space radius possess *some* structure corresponding, via the duality map, to a target space of *large* radius, which coincides with phenomenal spacetime), and since this structure does qualify as spatiotemporal on our proposed augmented Knoxian functional definition of spacetime, it seems that, ultimately, even on this revised functional definition, spacetime in perturbative string theory is non-fundamental but necessary after all.

8 Conclusions

A number of tasks have been accomplished in this paper. Having introduced **MR1** and **MR2** and their ‘miraculous’ status in GR in §2, and their connections with the chronogeometricity and spatiotemporality of the metric field in §3, I turned to perturbative string theory in §4. In that section, I demonstrated that **MR1** and **MR2** cease to be miraculous in this theory of quantum gravity—rather, they are necessary consequences of the formalism. This done, I considered the chronogeometricity and spatiotemporality of the target space metric field g_{ab} in perturbative string theory; since **MR1** and **MR2** hold of necessity in this theory, so too (on a Knoxian account) does the spatiotemporal status of g_{ab} . I clarified the logic of these arguments in §5.

In §6, I distinguished two senses of ‘emergent spacetime’: (A) the object which plays the spatiotemporal role in a given theory not being part of the fundamental ontology of that theory (as specified in its KPMs); versus (B) the object which plays the spatiotemporal role doing so in some, but not all, models of the theory. I argued that while (again on a Knoxian account) GR exemplifies (B) but not (A), perturbative string theory exemplifies (A) but not (B). Finally, in §7 I drew upon Huggett’s discussion of T-duality and phenomenal space to reflect upon whether these verdicts are reasonable, and whether both they, and Knox’s spacetime functionalism more generally, should be revised.

There are several important upshots of this work. First, (i) the coincidence of inertial and gravitational mass, or (ii) the geodesic principle, are arguably ‘miraculous’—

unexplained—in theories prior to GR, but find natural explanation in that context.⁴² The observation that **MR1** and **MR2** admit of natural explanation in perturbative string theory, where none was forthcoming in GR, can be seen as a continuation of this trend of accounting for apparent coincidences in one’s physics by recourse to a more ‘fundamental’ theory. (If one takes the explanation of apparent coincidences to be a theoretical virtue, then this is, of course, a reason to favour perturbative string theory over GR. I do not wish, however, to claim that these observations lend any epistemic support to the claim that perturbative string theory is *true*—for (a) it may be that there exist *other* candidate successor theories to GR which can also account for these coincidences [see below]; moreover (b) no explicit connection has yet been forged between a theory’s accounting for such coincidences, and its being true. For more on non-empirical theory confirmation in the context of string theory, see [15].) Second, this work is useful for clarifying senses of emergent spacetime, some of which apply to GR, and others to string theory; this illustrates that there is no univocal notion of the ‘emergence of spacetime’ in a given theory. Third, consideration of phenomenal spacetime in Huggett’s sense re-emphasises that more needs to be done to understand fully the spatiotemporal picture with which perturbative string theory presents us.

Possible extensions of investigations into the extent to which **MR1** and **MR2** admit of explanation in successor theories to GR abound. Essentially, what one needs in order to account for these ‘miracles’ is the existence of a unified origin for dynamical and metric symmetries. There are some hints that this can be attained not just in perturbative string theory, but also via ‘supergravity’ and ‘double field theory’.⁴³ The exploration of such matters should constitute a philosophically fertile research programme for the future.

⁴²See e.g. [7, 55] for discussion of such matters. For some discussion calling into question such claims regarding the geodesic principle in particular, see [32, 56, 57].

⁴³Many thanks, respectively, to Richard Dawid and Nic Teh for these suggestions. For introductions to these areas of theoretical physics, see respectively [37] and [1].

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