Flat Spacetime Gravitation with a Preferred Foliation

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Paralleling the formal derivation of general relativity as a flat spacetime theory, we introduce in addition a preferred temporal foliation. The physical interpretation of the formalism is considered in the context of 5-dimensional “parametrized” and 4-dimensional preferred frame contexts. In the former case, we suggest that our earlier proposal of unconcatenated parametrized physics requires that the dependence on \( \tau \) be rather slow. In the 4-dimensional case, we consider and tentatively reject several areas of physics that might require a preferred foliation, but find a need for one in the process (“flowing”) theory of time. We then suggest why such a foliation might reasonably be unobservable.

I. INTRODUCTION

It is known that general relativity can be formally derived as a flat spacetime theory[1,2]. Such a derivation can be based on invariance (up to a boundary term) of the free gravitational action under an infinitesimal gauge transformation and on universal coupling of all energy-momentum, including that of gravitation itself, to serve as the source[7,8]. In this context there exists a gravitational energy-momentum tensor[9], not merely a pseudotensor, so gravitational energy-momentum is localized in a coordinate-invariant (but gauge-variant) way. As W. Thirring observed, it is not clear a priori why Riemannian geometry is to be preferred over all the other sorts of geometry that exist, so a derivation (as opposed to postulation) is attractive[3].

Here we introduce on an \( n \)-dimensional space a preferred temporal foliation \( \partial_\mu \theta \) along side the \( n \)-dimensional flat metric \( \eta_{\mu\nu}(x) \) (in arbitrary coordinates) as nondynamical objects. While the dimension \( n \) is not specified, we expect that \( n = 4 \) with metric signature \(-+++\) and \( n = 5 \) with signature \(-++++-\) will be of the greatest interest. (If \( n = 4 \), then Greek indices run from 0 to 3 and \( \theta = x^0 = t \) for some preferred inertial frame in the natural coordinates. If \( n = 5 \), then Greek indices run over 0,1,2,3,5 with \( \theta = x^5 = \tau \) in natural coordinates, \( \tau \) being an invariant supertime.) The foliation \( \partial_\mu \theta \) obeys \((\partial_\mu \theta) \eta^{\mu\nu} \partial_\nu \theta = -1 \) and \( \partial_\mu \partial_\nu \theta = 0 \). We define the unit normal covector by \( n_\mu = -\partial_\mu \theta \). Its \( \eta \)-raised counterpart \( n^\nu = n_\mu \eta^{\mu\nu} \) is future-pointing[24]. The metric and foliation induce a projection tensor \( h_{\mu\nu} = \eta_{\mu\nu} + n_\mu n_\nu \). If we agree to refer to \( \theta \) as “time” (perhaps meaning \( \tau \)) and the remaining dimensions as “space” (perhaps having signature \(-+++\)) for convenience, then the projection tensor serves as the “spatial” metric. The symmetric gravitational potential is \( \gamma^{\mu\nu} \) (of density weight 0), and bosonic matter fields are denoted by \( u \) (with all indices suppressed). One could take \( \gamma \) to be either a contravariant density of any weight (except \( \frac{1}{2} \)) or a covariant density of any weight (except \( -\frac{1}{2} \)). (These weights are special because the resulting inverse metric tensor density or metric tensor density, respectively, has a determinant of \(-1\), so there are only 9 independent components. Thus, invertibility issues arise.) In the massless case, this choice makes no difference. If it were desired to add a gauge-symmetry-breaking mass term, the choice of index position and weight would affect the results.

II. SPATIAL GAUGE INVARIANCE AND UNIVERSAL MOMENTUM COUPLING

This derivation is an improved version of previous work of ours on this subject[23], based on improvements in the corresponding derivation of general relativity[1]. Unlike our previous parametrized derivation[23], this one does not assume that \( n_\mu n_\nu \gamma^{\mu\nu} = 0 \), so one can retain the lapse function \( N \) of an ADM split[21] as a nontrivial quantity, as opposed to requiring \( N = 1 \) a priori and thus not varying it. The assumption of gauge invariance requires that the field be massless, except for the time-time part \( n_\mu n_\nu \gamma^{\mu\nu} \), which can be massive. One can show this fact directly by writing the most general algebraic quantity that is quadratic in the gravitational potentials. (As in the case without a preferred foliation[23], there is no mathematical objection to adding a mass term that breaks the gauge invariance. A possible physical difficulty with negative energy will be noted below. Including both mass terms would imply that the time-time part of the field had a different rest mass from the rest of the field.)

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A. Free Field Action

Let $S_f$ be the action for a free gravitational field (in the sense that the gravitational coupling constant vanishes). We require that $S_f$ change only by a boundary term under the infinitesimal gauge transformation

$$\gamma_{\mu\nu} \rightarrow \gamma_{\mu\nu} + \partial^\mu \xi^\nu + \partial^\nu \xi^\mu,$$

(1)

$\xi^\nu$ being a vector field obeying $\xi^\nu n_\mu = 0$.

For any $S_f$ invariant in this sense under (1), the free field equations’ divergence is purely temporal, for its spatial projection vanishes, as we now show. The action changes by

$$\delta S_f = \int d^nx [\frac{\delta S_f}{\delta \gamma_{\mu\nu}} (\partial^\nu \xi^\mu + \partial^\mu \xi^\nu) + e^\nu,_{\mu}] = \int d^nx f^\mu,_{\nu}.$$

(2)

The forms of the boundary terms $e^\nu,_{\mu}$ and $f^\mu,_{\nu}$ are not needed. Integrating by parts, letting $\xi^\mu$ have compact support to annihilate the boundary terms (as we shall do throughout the paper), and recalling the purely spatial character of $\xi^\mu$, we obtain the identity

$$h^\nu,_{\alpha} \partial^\mu \frac{\delta S_f}{\delta \gamma_{\mu\nu}} = 0.$$

(3)

B. The Asymmetric Metric Energy-Momentum Tensor

An expression for the total energy-momentum tensor can be derived from $S$ using the metric recipe in the following way, making allowances for the presence of another absolute object. The action depends on the flat metric $\eta_{\mu\nu}$, the foliation $n_\mu$, the gravitational potential $\gamma_{\mu\nu}$, and bosonic matter fields $u$, the last representing any number of tensor densities of arbitrary weights and index positions. Under an arbitrary infinitesimal coordinate transformation described by a vector field $\psi^\mu$, the action changes by the amount

$$\delta S = \int d^nx \left[ \frac{\delta S}{\delta \gamma_{\mu\nu}} L_\psi \gamma_{\mu\nu} + \frac{\delta S}{\delta u} L_\psi u + \frac{\delta S}{\delta \eta_{\mu\nu}} L_\psi \eta_{\mu\nu} + \frac{\delta S}{\delta n_\mu} L_\psi n_\mu + g^\mu,_{\nu} \right],$$

(4)

where $L_\psi$ is the Lie derivative with respect to $\psi^\mu$ and $g^\mu,_{\nu}$ is another boundary term of no present interest. But $S$ is a scalar, so $\delta S = 0$. Letting the matter and gravitational field equations hold gives

$$\delta S = \int d^nx \left( \frac{\delta S}{\delta \eta_{\mu\nu}} L_\psi \eta_{\mu\nu} + \frac{\delta S}{\delta n_\mu} L_\psi n_\mu \right) = 0,$$

(5)

or, upon rewriting the Lie derivatives in terms of the covariant derivative compatible with the flat metric $L_\psi$, we have

$$\partial^\nu \frac{\delta S}{\delta \eta_{\mu\nu}} - \frac{1}{2} n_\nu \partial^\mu \frac{\delta S}{\delta n_\mu} = 0.$$

(6)

One can identify $\frac{\delta S}{\delta \eta_{\mu\nu}} - \frac{1}{2} n_\nu \eta_{\mu\nu} \frac{\delta S}{\delta n_\mu}$ as the energy-momentum tensor. Its asymmetry reflects the preferred character of $n_\mu$. If we take the spatial projection of this quantity, we obtain the tensor of momentum density and its flux $h^\nu,_{\alpha} \frac{\delta S}{\delta \eta_{\mu\nu}}$.

C. Full Universally-Coupled Action

For the full theory, we postulate that the spatial projection of the Euler-Lagrange equations should be just the spatial projection of the free field equations for $S_f$ augmented by the momentum tensor (including gravitational momentum):

$$h^\nu,_{\alpha} \frac{\delta S}{\delta \gamma_{\mu\nu}} = h^\nu,_{\alpha} \frac{\delta S_f}{\delta \gamma_{\mu\nu}} - \lambda h^\nu,_{\alpha} \frac{\delta S}{\delta \eta_{\mu\nu}},$$

(7)

where $\lambda$ is a coupling constant. If $n = 4$ and the theory in question is general relativity, then $\lambda = -\sqrt{32\pi G}$; for other theories, the Newtonian limit would need consideration. If $n = 5$, then $\lambda$ has different dimensions from Newton’s
constant $G$, with an additional length entering $\{22\}$. Horwitz et al. have previously noted the entrance of an additional length in their parametrized electromagnetism $[19]$. One is free to make a change of variables in $S$ from the flat metric $\eta^{\mu\nu}$ and gravitational potential $\gamma^{\mu\nu}$ to $\eta^{\mu\nu}$ and $g^{\mu\nu}$, where

$$g^{\mu\nu} = \eta^{\mu\nu} + \lambda \gamma^{\mu\nu}. \quad (8)$$

Equating coefficients of the variations gives

$$\frac{\delta S}{\delta \eta^{\mu\nu}}|_\gamma = \frac{\delta S}{\delta \eta^{\mu\nu}}|_g + \frac{\delta S}{\delta g^{\mu\nu}} \quad (9)$$

and

$$\frac{\delta S}{\delta \gamma^{\mu\nu}} = \lambda \frac{\delta S}{\delta g^{\mu\nu}}. \quad (10)$$

Putting these two results together gives

$$\lambda \frac{\delta S}{\delta \eta^{\mu\nu}}|_\gamma = \lambda \frac{\delta S}{\delta \eta^{\mu\nu}}|_g + \frac{\delta S}{\delta \gamma^{\mu\nu}} \quad (11)$$

Equation (11) splits the energy-momentum tensor into one piece that vanishes when the gravitational Euler-Lagrange equations hold and one piece that does not. Using this result in (7) gives

$$\lambda h^\nu_\alpha \frac{\delta S}{\delta \eta^{\mu\nu}}|_g = -h^\nu_\alpha \frac{\delta S_1}{\delta \gamma^{\mu\nu}}, \quad (12)$$

which says that the spatial projection of the free field Euler-Lagrange derivative must equal (up to a constant factor) that part of the momentum tensor that does not vanish when the gravitational field equations hold. Recalling (3), one derives

$$h^\nu_\alpha \partial^\mu \frac{\delta S}{\delta \eta^{\mu\nu}}|_g = 0, \quad (13)$$

which says that the part of the momentum tensor not proportional to the gravitational field equations has identically vanishing divergence. This result concerning the splitting of the momentum tensor will be used in considering the gauge transformations of the full theory. It also ensures that the gravitational field equations alone entail conservation of momentum, without any separate postulation of the matter equations. (Previously the derivation of a conserved momentum tensor required that gravity and matter obey their field equations, as in (5).) But the conservation of energy still depends upon the matter field equations.

Expanding the projection tensor gives

$$\partial^\mu \frac{\delta S}{\delta \eta^{\mu\alpha}}|_g + n^\nu n_\alpha \partial^\mu \frac{\delta S}{\delta \gamma^{\mu\nu}}|_g = 0. \quad (14)$$

Let us divide the action into 3 pieces: a piece $S_1$ independent of $\eta^{\mu\nu}$ (the other variable being $g^{\mu\nu}$, not $\gamma^{\mu\nu}$), a piece $S_2$ that contributes a symmetric curl (i.e., a quantity with identically vanishing divergence on either index) to the stress tensor, and another piece $S_3$. We therefore write

$$S = S_1[g, u, n] + S_2[g, u, \eta, n] + S_3[g, u, \eta, n], \quad (15)$$

suppressing all indices in the arguments. Given the assumed property of $S_2$, one can write

$$-\eta^{\mu\alpha} \eta^{\nu\beta} \frac{\delta S_2}{\delta \eta^{\mu\beta}}|_g = \frac{1}{2} \partial_\rho (M^{[\mu\rho]}|^{\nu}\sigma + M^{[\nu\rho]|^{\sigma}\mu}) + b \sqrt{-\eta} \eta^{\mu\nu} \quad (16)$$

[28] (pp. 89, 429), where $M^{\mu\nu\rho}$ is a tensor density of weight 1 and $b$ is a constant. This result follows from the converse of Poincaré’s lemma in flat spacetime. One easily verifies that if $3$

$$S_2 = \frac{1}{2} \int d^n x R_{\mu\nu\rho\sigma}(\eta) M^{\mu\nu\rho\sigma}(g, u, \eta, n) + \int d^n x \chi^\mu_{\nu, \mu} + 2b \int d^n x \sqrt{-\eta}, \quad (17)$$
then \( \delta S_3 / \delta g \) has just the desired form, while \( S_2 \) does not affect the Euler-Lagrange equations because its value is 0. \( \lambda^{\mu} \) is a weight 1 vector density, because we require that \( S \) be a scalar. For convenience we deposit all boundary terms into \( S_2 \).

One can also show that if

\[
S_3 = -\frac{1}{2} \int d^n x \partial_{\beta} \psi^{\alpha \beta}(g, u, \eta, n) n_\alpha \partial_{\mu} n^\mu,
\]

then (14) is satisfied. Now both terms of that equation are used, not just one as for \( S_2 \). Here \( \psi^{\alpha \beta}(g, u, \eta, n) \) is any (2,0) weight 1 tensor density constructed from all the fields, dynamical and nondynamical. It turns out that \( \delta S_3 / \delta g \) does not contribute nontrivially, one must recall from (13) above that

\[
\delta \psi^{\alpha \beta}(g, u, \eta, n) n_\alpha \partial_{\mu} n^\mu = 0.
\]

Combining these 3 terms gives the total action

\[
S = S_1[g^{\mu \nu}, u, n_\mu] + \frac{1}{2} \int d^n x R_{\mu \nu \rho \sigma}(\eta) M^{\mu \nu \rho \sigma}(g, u, \eta, n) + 2b \int d^n x \sqrt{|\eta|} +
\]

\[
-2 \int d^n x \partial_{\beta} \psi^{\alpha \beta}(g, u, \eta, n) n_\alpha \partial_{\mu} n^\mu + \int d^n x \lambda^{\mu}(g, u, \eta, n) n_\mu.
\]

One sees that \( S_1 \) can contain mass and self-interaction terms for the time-time component of the gravitational potential, using \( \sqrt{-g} \) to the first power and any suitable function of \( g^{\mu \nu} n_\mu n_\nu + 1 \). With this action, one could write down a number of theories, including (of course) general relativity. One would still need to verify that any given theory made sense theoretically (such as by having adequate positive energy properties) and empirically, for the universal coupling principle does not address such questions fully.

D. Gauge Invariance

It is instructive to determine what has become of the original free field gauge invariance. The scalar character of the action \( S[g, u, \eta, n] \) implies that the variation of the action under a coordinate transformations vanishes:

\[
\delta S_{\text{cf}} = \int d^n x (\delta g^{\mu \nu} / \delta \eta^{\mu \nu} \lambda^{\mu} g^{\nu} + \delta S / \delta u u + \delta S / \delta g^{\mu \nu} |g^{\nu} \eta^{\mu} + h^{\mu \eta} |) = 0.
\]

But in a flat spacetime theory, invariance under coordinate transformations is trivial. A gauge transformation, on the other hand, would be a transformation that changes the action changes only by a boundary term, but is not a coordinate transformation. Using the coordinate transformation formula and noting that the terms involving the absolute objects do not contribute more than a divergence, one easily verifies that a (pure) gauge transformation is given by \( \delta g^{\mu \nu} = \lambda \xi g^{\mu \nu}, \delta u = \xi u, \delta \eta^{\mu \nu} = 0, \delta n_\mu = 0, \) where \( \xi^{\mu} n_\mu = 0, \) but \( \xi^\mu \) is otherwise arbitrary. In showing that the term for the flat metric does not contribute nontrivially, one must recall from (13) above that

\[
\lambda^{\mu} \partial_{\mu} \frac{\delta S}{\delta \eta^{\mu \nu}} |g = 0
\]

identically. The term for the foliation does not contribute at all because \( \lambda \xi n_\mu = 0 \) on account of the constancy of \( n_\mu \) and the assumption that \( \xi^{\mu} n_\mu = 0 \). Thus, gauge transformations change (bosonic) dynamical fields in the same way that ‘spatial’ coordinate transformations do, but leave the nondynamical objects unchanged.

E. Bianchi Identities

Taking the independent variables to be those in \( S[g^{\mu \nu}, u, \eta^{\mu \nu}, n_\mu] \), one can easily derive the Bianchi identities. Letting the Euler-Lagrange equations hold, one finds an additional equation that holds as a consequence. This phenomenon is not unprecedented: in unimodular general relativity (which also has a nondynamical object present, namely, a volume element), this additional equation restores the trace of the Einstein equations (up to an arbitrary cosmological constant) that failed to appear in the Euler-Lagrange equations [29]. So a constraint not admitted through the front door might still reenter through the back door.
From a Hamiltonian point of view, one says that dynamical preservation of the momentum constraints (roughly, the time-space Einstein equations) implies the Hamiltonian constraint (roughly, the time-time component of the Einstein equations), up to an arbitrary constant, in the unimodular theory. On the other hand, we find that in the theory obtained from the Hamiltonian for general relativity but with \( N = 1 \) \textit{a priori} (so the Hamiltonian constraint does not follow from the variation of the Hamiltonian), no further constraint, including any portion of the Hamiltonian constraint, is required to preserve the momentum constraints. At least that is the case in the vacuum theory; we expect that minimally coupled matter would behave similarly. Also see (23). But it would appear that this theory would have trouble with the negative energy degree of freedom constituted by the determinant of the curved spatial metric, so it does not seem physically viable.

We return now to the Lagrangian formulation for theories with a flat background metric and a preferred foliation. Making a coordinate transformation, letting the matter and gravitational field equations hold, discarding the boundary terms, and using the arbitrariness of the coordinate transformation yields the relation

\[
\partial^\mu \frac{\delta S}{\delta y^\mu} |g - \frac{1}{2} n_\nu \partial_\mu \frac{\delta S}{\delta n_\nu} = 0. \tag{23}
\]

On account of (13), the spatial projection of this equation already holds identically. The remainder, the temporal component, takes the form

\[
2n_\nu \partial^\mu \frac{\delta S}{\delta y^\mu} |g + \partial_\mu \frac{\delta S}{\delta n_\mu} = 0. \tag{24}
\]

When the action is expanded as using \( S = S_1 + S_2 + S_3 \), as in (13), several terms vanish, namely, \( \frac{\delta S_1}{\delta y^\mu} |g, \partial^\mu \frac{\delta S_2}{\delta y^\mu} |g, \) and \( \frac{\delta S_3}{\delta y^\mu} \). After showing that \( \frac{\delta S_2}{\delta y^\mu} = 2 \partial^\mu n_\nu \partial_\mu \psi^\alpha \), one finds that the contributions from \( S_3 \) cancel each other. The final result takes the rather simple form

\[
\frac{\partial}{\partial x^\mu} \frac{\delta S_1}{\delta n_\mu} = 0. \tag{25}
\]

The physical meaning of this equation will depend on the precise form of \( S_1 \).

### III. INTERPRETATION OF THE \(--+-+--\) FORMALISM

In 5-dimensional form this work suggests itself as a route to a theory of “parametrized” gravitation, the extra dimension being the invariant supertime \( \tau \). If the extra time dimension is to be interesting, it is necessary that physical fields be permitted to depend on \( \tau \). But then one faces the question of relating a 5-dimensional description to the observed 4 dimensions. One common approach in an electromagnetic context has been “concatenation,” in which the \( \tau \)-dependent vector potential is integrated over all \( \tau \) (from eternity past to eternity future) to give standard Maxwell potentials, the latter supposedly being tied to experiments [19, 20, 11, 12, 14, 16, 17]. We have previously argued that concatenation is unsatisfactory [22], pace (23). One reason is that it makes essential use of the linearity of the field equations, but this linearity does not hold generally; in particular, it is violated by any reasonable theory of gravity [22]. It is doubtful that any variant of concatenation of a nonlinear parametrized theory would give a plausible nonlinear nonparametrized theory. Even Yang-Mills theories, though formally similar to electromagnetism, cannot be concatenated. Another reason is that observations are said to be influenced by all values of \( \tau \), including future ones, yet \( \tau \) is thought to be related to the process/flowing aspect of time [19]. The latter fact implies that experiments performed by real people in ordinary life ought to occur at some definite moment (or finite interval) of \( \tau \). So concatenation introduces the paradox of backwards causation in \( \tau \) [22]. It also introduces a curious distinction between measurement and evolution, like certain versions of quantum mechanics, rather than regarding measurement as a specific kind of evolution.

Motivated by these criticisms, we previously suggested that omitting concatenation and interpreting all experiences as involving a convective derivative with respect to \( \tau \) along a worldline might yield an adequate interpretation of the parametrized formalism [22]. However, the non-concatenated view has a drawback of its own, for such a theory generically agrees with standard well-confirmed 4-dimensional electromagnetism (e.g., light speed measurements) only if the dependence on \( \tau \) is quite weak [22]. This limit is similar to the zero mode limit considered by Frustai and Horwitz [21], who have been aware of some of the difficulties with concatenation [11]. They observed that the zero mode limit is a sufficient condition for agreement of the parametrized theory with experiment. But our suggestion
is that approaching the zero mode limit is a necessary condition. For example, the lack of observed dispersion in light propagation indicates that the $\tau$ derivatives are much smaller than the $t$ derivatives in such contexts. One could wonder what physical purpose $\tau$ serves [32]. If $\tau$ were to be associated with temporal becoming only, with no physical content (c.f. (33)), then perhaps a more attractive and economical solution exists in the $n = 4$ context, as we will explain below.

A recent development in membrane theory might possibly be of interest here. Recently the somewhat analogous question “Can there be ‘large’ extra dimensions?” has received a surprising positive answer. The Randall-Sundrum scenario [34,35] appears to permit higher-dimensional theories to have large extra dimensions while giving empirically reasonable results. Perhaps one can imagine a parametrized analog of this move.

Our discussion above has considered only parametrized field theory. It is also worthwhile to consider the interpretation of parametrized particle mechanics, a subject that has received some attention [18,36,19,37]. If the Stueckelberg-Feynman notion of positrons as electrons moving backwards-in-$t$ with respect to $\tau$ is employed, as has often been suggested, then the issue of backwards-in-$\tau$ causation appears. One might instead require that particle trajectories be such that $t$ is always an invertible function of $\tau$ to avoid that problem. One drawback is that one loses the opportunity to consider pair-creation and pair-annihilation processes in a classical framework. But then it appears that field theory is necessary after all to handle a variable number of particles, and the problems above must be faced.

In conclusion, it is clear that the parametrized formalism still faces fundamental interpretive questions.

We also note some interesting work by Cawley [38] and by Hsu and Shi [39]. The classical work of J. L. Cook [27,40] will be considered in an appendix.

IV. INTERPRETATION OF THE $-+-+-$ FORMALISM

In a 4-dimensional context this formalism corresponds to the existence of a preferred reference frame. It is generally assumed that no such thing exists, though the subject has received some attention [11,12]. The presumed nonexistence of a preferred frame is in some important respects quite helpful, because of the resulting tight restrictions on the number of theories that can be conceived. With a very few possible exceptions, all known physical processes are consistent with the orthodox relativistic view that there is no preferred foliation (and that backwards-in-time causation does not occur). In view of the apparently limited gains and substantial losses realized by giving up Lorentz invariance, one might wonder what is the purpose of considering a preferred foliation in physics. It would seem that a rather good argument is needed to justify such work. We now consider whether such an argument is available.

A. Possible Habitats for a Preferred Foliation in Physics?

One apparent difficulty for standard Lorentz-invariant physics is the remarkable behavior seen in certain quantum mechanical experiments, such as by Aspect et al., in which 2 particles in a suitable superposed state seem to be able to ‘communicate’ superluminally. Much of the physics community seems to believe that locality is doomed, and has given up on it, at least when its mind is on quantum mechanics. However, this would be a tremendous loss, and so it ought not to be accepted unnecessarily [43].

What can be said in defense of locality? We do not claim to give a comprehensive review, but only provide two suggestions. Evidently the detector efficiency loophole is still open [13,14]. Szabo and Fine’s model even works for experiments testing the GHZ scenario [14], a more recent and perhaps more potent threat to local hidden variables than Bell’s theorem. Detector efficiencies are still too low to close this loophole [44]. One also knows that the experiments violating the Bell inequalities are compatible with the orthodox relativity if one is prepared to embrace “superdeterminism” [14,49], which violates the inequalities by introducing correlations between the hidden variables and the detector settings. By positing a common cause for these correlations, one can preserve orthodox relativity, the Aspect experiments notwithstanding. Because the GHZ theorem involves similar locality assumptions to those involved in Bell’s theorem [50], we suspect that it can be subverted analogously. However, this view’s demanding philosophical underpinnings, such as its denial of (libertarian) free will[4] and evident need for an all-determining Agent

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1Free will faces a potent long-standing conceptual objection that an action that isn’t fully caused is to that extent merely
to correlate the initial conditions of the world, might limit its appeal (see, e.g., Bell’s attitude [51] (pp. 100-103, 110, 154)). The detector efficiency loophole has also seemed unappealing to some, such as Bell [51] (p. 109). However, it is at least worthwhile to show that these strategies exist, because they show that even in this peculiar aspect of quantum mechanics, nothing is presently known with certainty that requires a preferred frame.

Another trouble spot for the usual relativistic view of time is quantum gravity’s “problem of time” [56], which consists in the prima facie disappearance of time from quantum versions of general relativity. However, we suspect that the problem lies not in the lack of a particular preferred frame (a feature shared with special relativity; the success of standard field theory in other contexts suggests that this feature is not at fault), but in the lack of a preferred class of inertial frames peculiar to the form of gauge invariance of general relativity. That is, plausibly the ‘fault’ lies in how general relativity differs from special relativity (general covariance, i.e., lack of nondynamical objects), so one needn’t add structures unknown to special relativity to address the issue. One expects that the problem of time would disappear if one suitably introduced a nondynamical background metric into the equations of motion. Adding a small rest mass to the theory would be an obvious way to implement this procedure (and thereby obtain a nonvanishing Hamiltonian), if the traditional negative-energy objection to massive gravity [57] (appendix on “ghost” theories) can be overcome. M. Visser has recently suggested that it can [58]. One might also prefer that the curved metric respect the flat background’s null cone structure, a nontrivial condition that, to our knowledge, has not been successfully imposed in an attractive way.

Recently another suggested habitat for the violation of Lorentz invariance has appeared. T. Jacobson and D. Mattingly have suggested that there is “reason to doubt exact Lorentz invariance: it leads to divergences in quantum field theory associated with states of arbitrarily high energy and momentum. This problem can be cured with a short distance cutoff which, however, breaks Lorentz invariance” [60]. They then introduce an “aether” consisting of a dynamical unit timelike vector (or covector) field. Their aether, being dynamical and failing in general to define a preferred foliation (because the covector typically is not a gradient), differs from what we consider. Their divergence argument might give a good reason to consider a preferred foliation, such as we have considered here. But it seems premature to put too much reliance on this proposal.

B. A Preferred Foliation from the Process Theory of Time

If these areas of physics do not provide sufficiently strong evidence for the existence of an observable preferred foliation in physics, then one might ask if there are extra-physical reasons for considering a preferred temporal foliation in physics. It so happens that in the 20th century’s central debate in the philosophy of time [61], one of the two views, if established, would show that a preferred foliation exists at the most fundamental level. (Below we will find authors arguing that if a preferred foliation exists, then presumably it manifests itself in physics.) This is the debate about the objectivity or otherwise of temporal becoming, that is, the ‘flow’ of time [62,63]. Some physicists and philosophers, based to a large degree on the influence of relativity [63], incline toward the “block universe” view that regards all moments of time as ultimately equal in status; the notions of past, present, and future are regarded as illusory or mind-dependent. If the block view (also known as “stasis,” “B-theory,” or “tenseless”) is correct, then

random and thus un-free [52, 54], so the denial of free will might be inevitable on other grounds. If so, then the entrance fee for superdeterminism will decrease.

2On the other hand, the 3 major monotheistic traditions all have (or had) strands that affirm theological determinism: Pharisaic Judaism [53], Reformed/Calvinist Christianity, and Islam. That there might be a natural affinity here is suggested by the language (e.g., (47) about events being “already ‘written in a book’.” The resemblance to Psalm 139:16 (NASB) cannot be accidental:

Thine eyes have seen my unformed substance;
And in Thy book they were all written,
The days that were ordained for me,
When as yet there was not one of them.

3This issue was considered for (massless) general relativity by Penrose [59]. Massive theories differ from the massless case in several ways: (1) the gauge-dependence of the relation between the two null cones disappears with the gauge freedom; (2) the supposed difficulty with long-range divergence between the curved and flat null geodesics disappears, and (3) the relation between the two null cones will tend to depend on the specific mass term employed.
all facts can in principle be displayed on a (single) spacetime diagram. On the other hand, if a spacetime diagram (perhaps augmented by timeless mathematical and logical truths) cannot display all facts, because some facts have temporal properties inconsistent with such a representation, then the “process” (“A-theory,” “tensed”) view that affirms objective becoming will be established.

The process view of time receives some unexpected assistance from stasis advocate D. H. Mellor, who wrote \[63\] (pp. 4-5):

The tenseless camp often offers only weak inducements to join it: the relative simplicity of tenseless logic, for example, or its consonance with relativity’s unification of space and time. But tenseless time needs a stronger sales pitch than that. Tense is so striking an aspect of reality that only the most compelling argument justifies denying it: namely, that the tensed view of time is self-contradictory and so cannot be true.”

Mellor claims to find this needed contradiction in McTaggart’s paradox, but this claim is not generally accepted and indeed appears to be false [63]. If McTaggart’s paradox fails to demonstrate a contradiction, but Mellor’s judgement is otherwise to be trusted, then the process view of time wins already. But some will require a more compelling case, which we believe can be made.

There is an argument (largely from [68]; a more developed version is forthcoming [69,70]) that appears to disprove the block view by showing that it cannot accommodate certain facts [71,72]. In order to make one’s appointments on time, one frequently needs to know what time it is now. For example, if one wants to pay taxes to the American government in a timely way, one might want to know that “It is now April 2.” Otherwise, one might file many weeks or even years late, because one just would not know when to file [65]. This sort of fact, along with more general facts about what is occurring now, cannot be represented tenselessly [73,74], such as on a spacetime diagram, or known by a timeless being (even a divine one, as Kretzmann has in mind) i.e., one lacking temporal location and duration. Let us see why this is the case. On such a diagram, one might make a mark at “April 2” on the time axis, but this mark will soon be outdated, so it will no longer represent “now.” Trying to keep the “now” mark current would require continually erasing and drawing on the spacetime diagram, which is of course illegal, for one then has a succession of diagrams (a movie), not a single one. Neither will the fact “It is April 2 on April 2” be of any use, both because it is a tautology [61] and thus cannot inspire any action at all, and because it is always true [76] and thus cannot motivate action at any special moment. If one finds the use of a date label such as April 2 troubling (as if a substantivalist view of time might be to blame), one could substitute some ordinary occurrence: “A rooster is crowing now” would serve, provided that the rooster only crows shortly before tax day. So “now” points to one or more facts that the block universe cannot accommodate. We must therefore reject the following claim by H. Reichenbach [77] (pp. 16,17):

There is no other way to solve the problem of time than the way through physics . . . If time is objective the physicist must have discovered that fact, if [sic] there is Becoming the physicist must know it . . . It is a hopeless enterprise to search for the nature of time without studying physics. If there is a solution to the philosophical problem of time, it is written down in the equations of mathematical physics.

We must further differ from Reichenbach, who asserted that determinism would exclude becoming [77], for irreducibly tensed facts provide a ground for an objective flow of time, even if determinism is true. (Interestingly, Reichenbach concluded that physics in fact does ground time flow.) Thus, we conclude that a preferred foliation generated by the “moving now” exists.

C. Must a Preferred Foliation Be Physically Observable?

It seems natural to assume that if a preferred foliation exists, it ought to manifest itself in physics fairly readily [61]. Such an intuition indicates that Reichenbach’s claim, though too strong, was not wholly misguided. In contemplating

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4 If one is reluctant to make statements such as “2 + 2 = 4 in 1980,” one can admit a class of timeless truths also, but that seems unnecessary [64]. Facts such as “I am John Perry” [65], if they are nontrivial, might not fit on a spacetime diagram, but since their temporal properties are not the problem, we can set them aside.

5 In the positivist era of the 20th century, the question in the the title of this subsection might have received the answer “of course,” because the verificationist criterion for meaning would have said that it was meaningless to talk about entities that are unobservable in principle. But such thought-stopping replies need not detain us today.
the notion of “beables” for quantum field theory, J. S. Bell wrote:

As with relativity before Einstein, there is then a preferred frame in the formulation of the theory … but it is experimentally indistinguishable. It seems an eccentric way to make a world.  

6 (p. 180, ellipses in the original; see also p. 155). (This seems to have been Bell’s a priori judgement of the idea. While he thought it somewhat odd, he nevertheless thought it worthwhile to consider as “the cheapest resolution” of what he saw a posteriori to be a real difficulty posed by the Aspect experiments 49. ) Philosopher T. Maudlin, in attempting to make sense of quantum mechanics and reconcile it to relativity, suggests that backwards causation or a preferred reference frame might be the least unacceptable ways of doing so 70. Concerning the possibility of a preferred frame in making sense of quantum mechanics, Maudlin, perhaps having in view Einstein’s line that God is subtle but not malicious, writes:

One way or another, God has played us a nasty trick. The voice of Nature has always been faint, but in this case it speaks in riddles and mumbles as well. Quantum theory and Relativity seem not to directly contradict each other, but neither can they easily be reconciled. Something has to give: either Relativity or some foundational element of our world-picture must be modified . . . the real challenge falls to the theologians of physics, who must justify the ways of a deity who is, if not evil, at least extremely mischievous. 79 (p. 242)

So if one is persuaded that the flow of time is objective and that a preferred foliation ought to show itself readily in physics, then one might consider the formalism above for gravitation with \( n = 4 \), or perhaps some other way of including the foliation in physics. D. Bohm’s nonlocal deterministic version of quantum mechanics is perhaps presently the most vibrant work that assumes a preferred foliation of 4-dimensional spacetime; the theory is presently being applied to quantum gravity and quantum cosmology 80,81. (But its nonlocality is not easy to embrace, even if one can tolerate a preferred frame.) However, as Butterfield and Isham note, “[m]ost general relativists feel [that] this response is too radical to countenance: they regard foliation independence as an undeniable insight of relativity” 52,53.

We suggest that the following explanatory strategy might relieve this tension between the philosophical support for objective becoming and the dearth of physical support for a preferred foliation, at least if theism is plausible. Rather than regarding the inclusion of a physically invisible (or nearly so) preferred foliation as an “eccentric” 51 or “if not evil, at least extremely mischievous” 79 way to make a world, one might suggest that the Maker rendered the preferred foliation physically invisible as an act of benevolence to physicists. To put it more plainly, the world was made Lorentz-invariant to make physics easier 4. As we noted above, the requirement of Lorentz invariance so restricts the possible theories that the principle of Lorentz invariance answers a vast number of questions that would otherwise require laborious experimentation to settle. Writing down, e.g., all possible terms in linearized gravity, first given Lorentz invariance, and then given a preferred foliation, would give one a clear sense of the economy afforded by Lorentz invariance. (Perhaps other symmetries are amenable to a similar interpretation.) We will content ourselves with the simpler scalar and vector field cases. Viewed from a spacetime perspective, rendering the foliation invisible amounts to keeping the foliation from appearing nontrivially in the action, with only the flat metric present. Viewed from the perspective of space-at-a-time, it means that the foliation only appears in concert with the spatial metric, in such a way that, along with a fundamental constant with dimensions of velocity, neither the foliation nor the spatial metric appears alone, but only the two combined into an effective spacetime metric.

We will now write down all possible terms that could appear in the Lagrangian density for a real scalar field \( \phi \), restricting the equations of motion to be linear and to have (at most) second derivatives. The existence of fundamental constants with dimensions of velocity (\( c \)) and angular momentum (\( \hbar \)) will be assumed, although we choose units in which these constants have unit value. Because the equations of motion, which are perhaps more important than the Lagrangian density, are unchanged by the addition of a divergence, we will regard terms that are equal up to a divergence as equivalent to avoid overcounting, and will drop terms that are themselves divergences. For brevity, we use a vertical bar to denote the flat metric’s covariant derivative. We also the overdot notation to indicate the time derivative \( n^\mu \partial_\mu \). The possible terms lacking \( n_\mu \) are \( \phi_{\mu} \phi^{\mu} \) and \( \phi^2 \). The first term is so basic that we will assume that it must always be present, even if the foliation appears in the theory. The values of all other coefficients are relative to this one, which depends on the normalization only. (Our results thus cannot be criticized as inflated, being instead

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6 Another suggestion might be that the world was made Lorentz-invariant to make physics prettier.
perhaps a bit pessimistic.) If \( n_\mu \) is also present, then \( \dot{\phi}^2 \) is also available. Thus there are 2 unspecified constants in the case with a preferred foliation observable, but there is only 1 with it absent, in the case of a scalar field. This is a savings already, though a modest one. The vector case will be more compelling.

In the case of a vector field, the possible terms in the Lorentz-invariant case, taking into account the restrictions above are these: \( A^\mu A^\nu_{|\mu|} \), \( A^\mu A^\nu_{|\nu|} \), \( A^\mu A^\tau_{|\mu|} \), and \( \epsilon^{\mu\nu\alpha\beta} A_{\mu|\nu|} A_{\alpha|\beta|} \), leaving 3 unspecified coefficients. If Lorentz invariance is not required, there appear in addition \( A^\mu A^\nu n_\mu n_\nu \), \( A^\mu A^\nu n_\mu n_\nu \), \( A^\mu A^\nu n_\mu n_\nu \), \( A^\mu A^\nu n_\mu n_\nu \), \( A^\mu A^\nu n_\mu n_\nu \), and \( \epsilon^{\mu\nu\alpha\beta} A_{\mu|\nu|} A_{\alpha|\beta|} \), giving 10 unspecified coefficients without Lorentz invariance. The economy afforded by Lorentz invariance is thus considerable in the case of a vector field. One expects that it would be substantial for higher-rank fields as well. We have not considered complex fields or fermionic matter, but we imagine that Lorentz invariance provides a respectable simplification in those cases as well. Complex fields have an interesting property not shared by a single real scalar field, viz., they admit first-order-in-time equations of motion such as the Schrödinger equation. (For a single real field, the most similar term in the Lagrangian density, \( \phi \dot{\phi} \), is merely a divergence.) So the investigation of these other types of fields, and of sets of real fields with internal symmetry groups, might be of interest.

In view of these simplifications, it is not too implausible to think that temporal becoming is objective, and yet physics is exactly Lorentz invariant, if the existence of a benevolent God who supports the enterprise of physics is plausible. Such a divine motivation does fit naturally within traditional monotheistic religion. According to the traditional story (Genesis 1:28), God tells the human race to reproduce and to fill the earth and subdue it, and to rule over other creatures. This command, which has been dubbed the “cultural mandate,” has been seen as encouraging the scientific study of natural phenomena. If this conclusion is correct, then physicists and philosophers can pursue their respective visions of time without mutual interference. The great gulf that Stapp [84] and Horwitz et al. [19] have found between “Einstein time” and “process time” is thus perhaps not so fixed. The need for a preferred foliation in physics would then need to be shown on grounds other than the flow of time, if at all.

This suggestion just made also answers an objection against the idea that God is temporal, i.e., has a location and a duration in time [27–88]. A temporal God’s knowledge of which events are objectively simultaneous defines a preferred foliation. But we have already addressed the relativistic objection to a preferred foliation. Introducing an omnipresent knower creates no further difficulties.

We close by concluding that even if one is persuaded of the existence of a preferred foliation of spacetime, there presently seems to be no compelling reason for rejecting standard Lorentz-invariant physics. However, there are some reasons strong enough to make consideration of the violation of Lorentz invariance an interesting pursuit.

APPENDIX A: \( \tau \)-RELATED WORK OF J. L. COOK

The work by J. L. Cook [27,40] deserves examination. In view of its early date for \( \tau \)-related work and its degree of development (to the point of including numerical results), this work perhaps deserves a larger role in the history of parametrized physics than it has hitherto received, if it is sound. However, we will raise questions about both the particle mechanics and the derivation of the 2-body gravitation interaction from a field theory.

Concerning mechanics, there is a question about how is the constraint of unit 4-velocity is preserved. Such a constraint is never mentioned by Cook, and it is suggested that particles have 4 degrees of freedom [27] (p. 122) [40] (p. 471). But given that Cook regards \( \tau \) as proper time [27] (p. 121), this constraint ought to hold. It is an important feature of relativistic mechanics, so its omission is puzzling.

In addressing the 2-body problem, Cook undertakes to derive the 2-body electromagnetic and gravitational interactions from corresponding field theories. For electromagnetism, \( \tau \)-dependent potentials are introduced in some postulated field equations. The justification for the equations seems to be not so much a principled \( \tau \)-dependent theory (which might follow from an action principle), but a desire to obtain equations that, upon integration over all \( \tau \) (a procedure later called “concatenation” by Horwitz et al., and which Cook sees as implicit in the work of Wheeler and Feynman), yield Maxwell’s equations. While an extra scalar potential in some way corresponding to \( \tau \) is introduced, it seems to play little role in the theory.

In view of the similarity between electromagnetism and general relativity, one might expect that the search for the gravitational 2-body interaction would follow the same plan as in electromagnetism. Surprisingly enough, that is not the case. Study of Cook’s work leaves us with significant unanswered questions.

Regarding the field theories used to derive the 2-body potentials, why is it that for gravitation, 5-dimensional
\(\tau\)-dependent Einstein field equations are assumed to hold, but 5-dimensional Maxwell equations are not assumed for electromagnetism \[27\]. This inconsistency suggests a degree of arbitrariness in the assumed field theories.

It should also be noted that the 5-dimensional Einstein equations are written in terms of 2-body relative coordinates, the center of mass motion presumably having been factored out. What does it mean to write a field theory in terms of relative coordinates? The relative coordinates in a 2-body problem take only 1 value at any moment (meaning a moment of \(\tau\), in this case), but the coordinates in a field theory take on all values. Furthermore, on what grounds does one assume that something like 5-dimensional Einstein field equations should hold with respect to the relative coordinates? This is hardly follows from “standard methods” (p. 127) or Møller’s book \[89\], to which Cook appeals. The situation becomes still less clear for the \(n\)-body problem, for then there are \(n-1\) sets of coordinates available.

Finally, concatenation of the linearized gravitational potential plays an essential role in finding the line element \[27\]. As we have noted above, concatenation of a \(\tau\)-local theory to give an ordinary theory depends essentially on the linearity of the theory \[22\], but linearity holds only in a weak-field approximation for gravitation. What is the generalization to the case of the full nonlinear Einstein equations that are assumed? Cook gives no hint of an answer, and we cannot think of any attractive candidates either.

In view of the questions listed above, we conclude that, although “interactions over surfaces of equal proper times are not known and must be determined from field equations” \[27\], the gravitational interaction found in \((27, 40)\) has not been persuasively derived from a field theory.

APPENDIX: ACKNOWLEDGMENTS

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\[\text{We now turn to amend a previous criticism of Cook’s paper to the effect that the allowed dependence on } \tau \text{ was restricted. This remark was based on the lack of } \tau \text{-derivatives in equation 46 of (22), based on the definition of the } \Box \text{ symbol in equation 42b. However, from equation 44 and the coordinate condition just above equation 46, it follows that the } \Box \text{ symbol in equation 46 ought to refer to a 5-dimensional wave equation, not a 4-dimensional one. One could then regard the lack of } \frac{\partial^2}{\partial \tau^2} \text{ in equation 46 as merely a typographical error.}\]
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