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INTRODUCTION

Despite the impressive array of techniques we currently possess for constructing initial data in general relativity, an unsatisfying feature persists: the sensitivity of the solution to the way we foliate spacetime. Indeed, this often overshadows entirely the physics we are attempting to understand. It is clear that the problem is even worse at the dynamical level.

In order to solve the constraints it is traditional to appeal to an extrinsic time foliation of spacetime. This involves some restriction on the extrinsic curvature so as to define the slicing. One such foliation, in particular, has a long history: that the spatial geometry be flat. It is well known that such slices are provided in the Schwarzschild spacetime geometry by Lemaître coordinates (see, e.g., [6]). Recently, these coordinates have been exploited to describe the canonical reduction of the theory with a source consisting of a shell of massive dust [7]. Indeed, there has been a revival of interest in the canonical reduction of spherically symmetric general relativity (see, for example, [8,9] and references therein). However, when the sources are not confined to a shell, the canonical procedure in an extrinsic time foliation is not generally tractable. The notable exception is the polar slicing [10,11]. But simplicity comes at a price: the corresponding chart is pathological if the slice possesses an apparent horizon, as it ultimately will in a classically collapsing geometry. Flat slicings, while remarkably simple, escape this shortcoming.

Historically, flat foliations of spacetime are of interest due to the role they played in the seventies as a provider of useful guidelines for the construction of the proof of the positive mass theorem. Brill and Jang [12] showed that if the spatial geometry is flat, and the extrinsic curvature falls off more rapidly than $r^{-3/2}$ at infinity, the only solution of the vacuum Einstein equations is flat. One might, therefore, be tempted to conclude that such foliations are pathological, inconsistent as they appear to be with the presence of gravitational waves. As we will see, however, the falloff of the extrinsic curvature in this gauge is slow (exactly as $r^{-3/2}$) and thus Brill and Jang’s argument is not applicable. Indeed, this slicing possesses the unusual feature that the Arnowitt-Deser-Misner (ADM) mass is encoded completely by the extrinsic curvature. While we do not have gravitational waves to contend with in a spherically symmetric geometry, it is still possible that the gauge exhibit pathologies associated with the
sources. However, we demonstrate here that if the sources
are finite, all solutions of the constraints are non-singular if
they are non singular at the origin. Therefore, this gauge
steers clear of singularities because the extrinsic curvature
(which is now the only measure of the geometry) is always
bounded when the sources are. This should be contrasted
with the situation in an extrinsic time foliation where the
solutions may well be singular if the sources are large but
finite [2,3].

We examine the occurrence of apparent horizons in the
intrinsic time gauge. We must abandon the geometrically
satisfying but nonetheless gauge dependent identification of
an apparent horizon with the formation of a bag with a neck
(minimal surface) in the spatial geometry. Because the ge-
ometry is flat, there is no bag to speak of, much less the kind
of bag envisaged by Wheeler. Now apparent horizons occur
due to the action of extrinsic curvature alone. However, it is
no longer enough that the extrinsic curvature be large, it also
must possess the appropriate sign. This is due to the pecu-
nularity of this gauge that, when the dominant energy condi-
tion is satisfied, the tangential projection \( K_R \) of the extrinsic
curvature tensor has a definite sign.

We next search for necessary and sufficient conditions for
the appearance of apparent horizons. Remarkably, once we
take into account the obvious obstruction associated with
the sign of \( K_R \), we can reproduce the inequalities we derived
earlier in the extrinsic time gauges [4,5]. This suggests a
validity which transcends their gauge dependent derivation.

It remains to confirm that the flat foliation is consistent
with the dynamics. We do this by constructing explicitly
both the lapse and the shift necessary to maintain the flatness
of the intrinsic geometry and we write down explicitly the
exterior form of the spacetime metric.

Schoen and Yau in [19] gave a sufficient condition for the
appearance of trapped surfaces in general initial data sets. It
turns out that the only place these estimates can be realisti-
cally tested is on intrinsically flat manifolds. We compare the
Schoen and Yau sufficient condition with ours for the case of
a constant density intrinsically flat sphere and find that ours
is better by a factor of 3.

Finally, in the same spirit which has prompted us to con-
sider flat foliations in the first place, we examine the flat
foliation as a special case of foliations with negative spatial
scalar curvature.

**FLAT FOLIATIONS**

We recall that the constraints are given by (we exploit the
notation introduced in [1])

\[
K_R[K_R+2K_\mathcal{L}] - \frac{1}{R^2} [2(2R')' - R'^2 - 1] = 8\pi\rho,
\]

\( (1) \)

and

\[
K'_R + \frac{R'}{R} (K_R - K_\mathcal{L}) = 4\pi J,
\]

\( (2) \)

where the line element of the spherically symmetric spatial
geometry is parametrized by

\[
ds^2 = dt^2 + R^2 d\Omega^2,
\]

and we have expanded the extrinsic curvature \( (n^a) \) is the out-
ward pointing unit normal to the two-sphere of fixed \( l \) as,

\[
K_{ab} = n_a n_b K_\mathcal{L} + (g_{ab} - n_a n_b) K_R.
\]

\( (4) \)

All derivatives are with respect to the proper radius of the
spherical geometry, \( l \). The spatial geometries we wish to
consider are \( R^3 \) \((l \in 0, \infty)\) with a single asymptotically flat
region with a regular center, \( l = 0 \). The appropriate boundary
condition on the metric at \( l = 0 \) is then

\[
R(0) = 0.
\]

\( (5) \)

We recall that \( R'(0) = 1 \) if the geometry is regular at this
point. We assume that both the energy density of matter, \( \rho \),
and its radial flow, \( J \), are appropriately bounded functions of
\( l \) on some compact support.

When we consider general spherically symmetric initial
data we start off with six functions, \( (g_{rr}, g_{\theta\theta}, K_R, K_\mathcal{L}, \rho, J) \),
which satisfy the two constraints. Because the dynamics re-
side in the matter field, so that the geometry is purely kine-
matical, it seems natural to choose \((\rho, J)\) as the independent
variables. This still leaves us with four dependent objects
satisfying the two constraints. One of the extra degrees of
freedom is obviously the coordinate choice on the three slice.
We can fix this more-or-less independently of everything
else. One natural choice (which we use) is to set \( g_{rr} = 1 \); the
other standard choices are to arrange that the metric be con-
formally flat or that the radial coordinate be the areal
(Schwarzschild) radius.

This leaves us with one extra variable among the three
\((g_{\theta\theta}, R^2, K_R, K_\mathcal{L})\) so we choose some relationship between
them. Such a condition should fix the slicing, the way that
the given slice is embedded into the spacetime. When we
solve the constraints completely we determine both the in-
trinsic geometry and the extrinsic curvature.

In this paper we propose to foliate spacetime using the
intrinsic geometry to mark time. The simplest possible
choice is the one we will adopt. Let us suppose that

\[
R(l) = l
\]

\( (6) \)

everywhere so that the spatial geometry is flat \( R^3 \) every-
where. This condition is, a priori, no more restrictive than
any of the other slicing conditions we have considered. The
scalar curvature now vanishes with the result that the Hamil-
tonian constraint reduces to the algebraic condition on \( K_{ab} \) in
terms of \( \rho \),

\[
K_R[K_R+2K_\mathcal{L}] = 8\pi\rho.
\]

\( (7) \)

The momentum constraint, apparently at least, is only modi-
ﬁed in a trivial way: \( R'/R \) is replaced by \( 1/l \).

To solve the constraints we eliminate the extrinsic curva-
ture scalar \( K_\mathcal{L} \) from Eq. (7) in favor of \( K_R \) and \( \rho \) and substi-
tute into (2). We get
\[(i^{3/2} K_R)' = 4 \pi i^{3/2} \left( \frac{\rho}{l K_R} + J \right). \tag{8}\]

The equation is no longer linear in \(K_R\). This is not too surprising. The constraints in general relativity are non-linear to start with and the physics is now completely encoded in the extrinsic geometry.

**Lemaître slicing of the Schwarzschild geometry**

As mentioned above, one example of such a flat slicing is given by the Lemaître coordinatization of the Schwarzschild solution. We start off with the Schwarzschild solution in standard coordinates

\[ds^2 = - \left(1 - \frac{2m}{R}\right) dt^2 + \left(1 - \frac{2m}{R}\right)^{-1} dR^2 + R^2 d\Omega^2, \tag{9}\]

and change the time coordinate via

\[t = \tau + 4m \sqrt{\frac{R}{2m} + 2m \ln \left| \frac{\sqrt{R} - \sqrt{2m}}{\sqrt{R} + \sqrt{2m}} \right|}. \tag{10}\]

Differentiating Eq. (10), we get

\[dt = d\tau + \left(1 - \frac{2m}{R}\right)^{-1} \frac{\sqrt{2m}}{R} dR, \tag{11}\]

and when this is substituted into Eq. (9) we obtain

\[ds^2 = - \left(1 - \frac{2m}{R}\right) d\tau^2 - 2 \sqrt{\frac{2m}{R}} R d\tau + dR^2 + R^2 d\Omega^2. \tag{12}\]

It is clear from Eq. (12) that the spatial geometry of the \(\tau\) equal constant slices is flat [14]. In addition, Lemaître coordinates are non-singular on the horizon at \(R = 2m\). The price one pays, however, is that the form of the spacetime metric is no longer diagonal in \(R\) and \(\tau\). Lemaître coordinates cover half of the maximal extension of the Schwarzschild geometry. For our choice, this is the region above the principal diagonal on the Kruskal diagram. Note that the spatial geometry is regular at the Schwarzschild singularity at \(R = 0\). Each spatial slice intersects this singularity.

It is simple to construct the extrinsic curvature tensor. The Wald ([13]) definition of the extrinsic curvature gives

\[\partial_0 g_{ab} = 2 N K_{ab} + N_{a;b} + N_{b;a}. \tag{13}\]

In the coordinate choice given by Eq. (12) we have \(\partial_0 g_{ab} = 0\), \(N = 1\), and \(N_i = \left( - \sqrt{2m/IR}, 0, 0 \right) \) which gives

\[K_{RR} = - \frac{1}{2} \sqrt{\frac{2m}{R^3}}, \quad K_{\theta\theta} = + \sqrt{2mR}, \]

\[K_{\phi\phi} = + \sqrt{2mR}\sin^2 \theta. \tag{14}\]

From the definition (4) we get

\[K_{\ell\ell} = \frac{3}{2} \sqrt{\frac{2m}{R^3}}, \quad K_{\ell\phi} = 0.\]

Clearly these satisfy Eqs. (7) and (8) with vanishing sources. Both \(K_L\) and \(K_R\) diverge at \(R = 0\). The singularity in the spacetime curvature is encoded completely in the extrinsic curvature.

**Momentarily static solutions**

A peculiarity of the flat foliation is that \(K_R\) responds directly to \(\rho\). Unlike its behavior in extrinsic time foliations, \(K_R\) does not vanish if \(J\) is identically zero unless \(\rho\) also is identically zero. Clearly, if there are any sources acting, this foliation does not admit any moment of time symmetry solutions with \(K_{ab} = 0\). In this respect, the simple constant density “star” solutions, with \(\rho\) constant on some compact support and \(J = 0\), should not be confused with genuine constant density solutions occurring in a momentarily static configuration.

**Two solutions for each \((\rho, J)\)**

One can show using Eq. (8) that there are two solutions for each specification of \((\rho, J)\). We write Eq. (8) as

\[K_R' + \frac{3}{2} l^{-1} K_R = 4 \pi l^{-1} \frac{\rho}{K_R} + 4 \pi J. \tag{16}\]

Suppose that \(K_R\) is finite at the origin, \(l = 0\). Then

\[K_R(0) = \pm \sqrt{\frac{8 \pi (0)}{3}}, \tag{17}\]

which is independent of \(J\). Having fixed \(K_R(0)\), Eq. (16) allows us to integrate out to find \(K_R(l)\).

**The momentum constraint and the quasi-potential local mass**

An alternative casting of the momentum constraint is as

\[(i^3 K_R')' = 8 \pi i^3 (\rho + JlK_R). \tag{18}\]

We recall that in this gauge the spherical quasi-local mass \(m\) (the Misner-Sharp, Hawking, et al. mass) [15,16] (see also [17] and [1] for references), defined by

\[m(l) = \frac{1}{2} R (1 - R^2) + R^3 K_R^2, \tag{19}\]

is completely determined by extrinsic curvature:

\[m(l) = \frac{1}{2} l^3 K_R^2. \tag{20}\]

and our rewriting of the constraint, Eq. (18), can be identified with the integrability condition on \(m\),

\[m' = 4 \pi R^2 (\rho R' + J R K_R). \tag{21}\]
Thus we can solve Eq. (18) as

$$K_R = \pm \frac{1}{\sqrt{2m(l)}} \frac{1}{l^3}, \quad (22)$$

where $m(l)$ is the quasi-local mass. Therefore, a necessary condition for the existence of a flat slice through a given spherical spacetime is that the mass function be positive. $m$ is manifestly positive in this gauge.

The geometry is necessarily singular if the constant of integration $m(0)$ is nonvanishing. Indeed, in vacuum, $m(l) = m(0)$ and we reproduce a Lemaître slice of the Schwarzschild geometry.

Outside the support of matter,

$$K_R^2 = \frac{2m(l_0)}{l^3}. \quad (23)$$

The ADM mass is encoded completely in $|K_R|$. We note also that $K_R$ has the same asymptotic falloff as the gauge, $K_L = 0.5K_R = 0$. In this gauge, the Hamiltonian constraint assumes the intrinsic geometric form,

$$\frac{3}{2}R = 16\pi\rho, \quad (24)$$

where $3R$ is the scalar curvature. It is equally legitimate, however, to treat Eq. (24) itself as a source dependent intrinsic time gauge condition. The Hamiltonian constraint then reduces to the algebraic condition, $K_L + 0.5K_R = 0$. This way, the extrinsic time gauge can be viewed as entirely equivalent to an intrinsic time gauge, albeit an unconventional one depending explicitly on the source.

**J = 0 is exactly solvable**

When $J = 0$, for any given $\rho$, Eq. (18) is exactly solvable. Suppose that the weak energy condition, $\rho \geq 0$, holds. Then

$$K_R = \pm \frac{1}{\sqrt{2M(l)}} \frac{1}{l^3}, \quad (25)$$

where [18]

$$M(l) = 4\pi \int_0^l d\ell\ell^2\rho, \quad (26)$$

is the bare material mass and equals the mass function. It must be positive. One solution is minus the other and $K_R$ has a definite sign. Indeed, when the dominant energy condition holds, this property continues to hold.

**$K_R$ has a definite sign**

If $\rho \geq c|J|$ for any constant $c$, then $K_R$ possess a definite sign. In particular, this is true if the dominant energy condition, $c = 1$, is satisfied. To prove this, let us suppose that $K_R$ is positive at $l = 0$. Suppose that it falls to zero on the support of $\rho$. Where $K_R > 0$, the combination $\rho/lK_R + J$ appearing on the right-hand side (RHS) of Eq. (8) is positive so that the LHS, $(1/2K_R)'$, is also positive. $K_R$ is therefore increasing and cannot fall through zero. An identical argument applies for negative $K_R$. In this gauge, dominant energy places a very strong constraint on $K_R$. Note, however, that the radial extrinsic curvature, $K_L$, which is now determined by Eq. (7), does not possess a definite sign.

The effect of introducing a current source, $J$, on a solution with $J = 0$ is easily deduced from Eq. (18). If $K_R$ is positive, then a positive current increases $K_R$, whereas a negative one decreases it. In particular, if $J$ has a definite sign (positive say), the solution has everywhere greater $K_R$ than the corresponding solution with $J = 0$.

**The extrinsic geometry is non-singular everywhere**

When $\rho$ and $J$ are finite, and the center is non-singular, $K_R$ is finite and bounded away from zero everywhere. However, the only way the extrinsic geometry can become singular is by $K_R$ diverging or $K_L$ diverging (for which it is necessary that $K_R = 0$) so that the flat slice is singularity avoiding.

We can place an explicit bound on $K_R$ as follows. We integrate Eq. (18)

$$L^2K_R^2 = 8\pi \int_0^L d\ell\ell^2(\rho + JK_R). \quad (27)$$

We then have

$$L^2(K_R)^2_{\text{Max}} \leq 8\pi \left[ \frac{1}{3}L^2\rho_{\text{Max}} + \frac{1}{4}L^3J_{\text{Max}}(K_R)_{\text{Max}} \right]. \quad (28)$$

so that (by solving the quadratic)

$$L(K_R)^2_{\text{Max}} \leq \pi J_{\text{Max}}L^2 + \left[ \pi^2(J_{\text{Max}}L^2)^2 + \frac{8\pi}{3}\rho_{\text{Max}}L^2 \right]^{1/2}. \quad (29)$$

$K_R$ is clearly finite when $\rho$ and $J$ are. We note that the bound is more sensitively dependent on $J_{\text{Max}}$ than it is on $\rho_{\text{Max}}$. Finally, we have seen that $K_R$ cannot vanish on the support of $\rho$ when the (generalized) dominant energy condition is satisfied. Thus $K_L$, which is determined by Eq. (7), is also bounded.

**APPARENT HORIZONS**

We define the two optical scalars,

$$\omega_\pm = 2(R^\perp \pm RK_R). \quad (30)$$

A future (past) apparent horizon forms when $\omega_\pm = 0$. The surface is future (past) trapped when $\omega_\pm \leq 0$. In this gauge, the appearance of an apparent horizon is entirely due to the action of extrinsic curvature.

We have already seen that solutions fall into two categories when the dominant energy condition is satisfied: those with $K_R > 0$ and those with $K_R < 0$. On solutions with $K_R$ positive (negative), future (past) trapped surfaces are impos-
We can exploit the universal bound on $K_R$, Eq. (29) to obtain a corresponding necessary condition.

Let the first apparent horizon occur at $l = L$. Then

$$L = 8\pi \int_0^L dl l^2 (\rho + JlK_R),$$

so that

$$L \leq \frac{8\pi L^3}{3} \rho_{\text{Max}} + 2\pi L^4 J_{\text{Max}}(K_R)_{\text{Max}}.$$  

We now exploit the bound Eq. (29) for $L(K_R)_{\text{Max}}$ to obtain

$$1 \leq 4\pi^2 (J_{\text{Max}}L^2)^2 + \frac{3}{\pi} \rho_{\text{Max}}L^2.$$  

On rearrangement, Eq. (39) can be cast

$$1 \leq 4\pi^2 (J_{\text{Max}}L^2)^2 + \frac{3}{\pi} \rho_{\text{Max}}L^2.$$  

If, in addition, we exploit dominant energy, we can replace $J_{\text{Max}}$ by $\rho_{\text{Max}}$ and solve the quadratic to get

$$\frac{1}{3\pi} \left( \frac{\sqrt{13}}{2} - 1 \right) \leq \rho_{\text{Max}}L^2.$$  

Thus if $\rho_{\text{Max}}L^2 < (\sqrt{13}/2 - 1)/3\pi$, the region cannot contain an apparent horizon. The constant of proportionality is comparable to that appearing in Eq. (57) of [5] with $\alpha = 1$. The derivation is, however, considerably simpler. In the extrinsic time foliation, the derivation depended in an essential way on the application of a weighted Poincaré inequality. This improvement is clearly related to the singularity avoidance of the intrinsic time gauge we have exploited here.

**CONSISTENCY OF FLAT FOLIATION WITH THE EVOLUTION**

The condition that the flat foliation be preserved under evolution, $\partial_0 \gamma_{ab} = 0$, implies that the extrinsic curvature is proportional to a Killing form:

$$2NK_{ab} = -\nabla_a N_b - \nabla_b N_a.$$  

In the spherically symmetric geometry we are considering this reduces to the set of conditions

$$N_i = NK_L$$

$$N_i = lNK_R.$$  

Here $N$ and $N_i$ are respectively the lapse and the radial shift. Eliminating $N$, we obtain
with the solution,

\[ N_l = \exp \left( - \int \frac{dl}{a(l)} \right) \]

\[ N = \exp \left( - \int \frac{dl}{a(l)} \right) / l K_R, \]  

where \( \alpha = -K_E/K_R \) is the ratio of extrinsic curvature scalars introduced in [3]. Thus flat slicing is consistent with evolution. Indeed, both the lapse and shift are completely determined without appealing to the dynamical Einstein equations for \( K_E \) and \( K_R \).

In the exterior region we have \( \alpha = 1/2 \), so that \( N_l = N_f(L(l))^{1/2} \) and \( N = N_f(L(l))^{1/2} / l^{1/2} \) where \( N_f(L) \) is the boundary shift. The lapse is constant. The boundary condition \( N \to 1 \) at infinity therefore fixes \( N_f(L) \). Let \( \tau \) be the time coordinate defined by this foliation. We then have the exterior spacetime metric

\[ ds^2 = - \left( 1 - \frac{2m_0}{l} \right) d\tau^2 \pm 2 \sqrt{\frac{2m_0}{l}} d\tau dl + dl^2 + l^2 d\Omega^2. \]

This is the Schwarzschild metric expressed in Lemaître coordinates as we have already seen in Eq. (12). The spacetime is completely characterized by the shift.

**NEGATIVE SCALAR CURVATURE**

Let us consider instead of the flat slicing, any foliation with negative scalar curvature:

\[ \sqrt{3} \mathcal{R} = - \frac{1}{K^2} \left[ 2 (R')^2 - R^2 - 1 \right]. \]

It is then simple to demonstrate that \( R' \geq 1 \) everywhere so that \( R \geq l \). We have, at a critical point of \( R' \),

\[ R'^2 = 1 - \sqrt{3} \mathcal{R} R^2. \]

If the scalar curvature is bounded then so is \( R' \). Such foliations are clearly very different from the extrinsic time foliations we considered in [3–5] with everywhere positive scalar curvature in which \( R'^2 \leq 1 \) and \( R \leq l \) in regular solutions. With a prescribed value of the scalar curvature, we can solve the constraints exactly as we did in the flat slice. We get

\[ (R'^2 K_R)^3 = 4 \pi R^{3/2} \left( \frac{R' \tilde{\rho}}{R K_R} + J \right), \]

where we set \( \tilde{\rho} = \rho - 3 \mathcal{R}/16 \pi \). We have \( \tilde{\rho} \geq \rho \) so that whatever energy condition is good with \( \rho \) is better with \( \tilde{\rho} \). Now, exactly as in a flat slicing, \( K_R \) has a definite sign when dominant energy holds.

It is straightforward to demonstrate that a bound can be placed on \( K_R \). Such gauges therefore share with the flat foliation its singularity avoidance. This differs from any foliation with \( \sqrt{3} \mathcal{R} \) positive where large sources are not always consistent with any singular geometry.

Let us examine the robustness of Eq. (35). It is simple to demonstrate that instead of Eq. (35), we obtain the sharp inequality

\[ RR' \geq M - \mathcal{P}, \]

which does not depend explicitly on \( \sqrt{3} \mathcal{R} \). Equation (50) coincides with Eq. (35) when \( R = l \). We see, however that the more negative the scalar curvature, the greater is the maximum of \( R' \), and thus the weaker the inequality.

**THE SCHÖEN AND YAU CRITERION FOR TRAPPED SURFACES**

The first mathematically precise statement of a sufficient condition for the appearance of apparent horizons in a general (i.e., nonsymmetric) initial data set was given by Schoen and Yau in [19]. This condition is difficult to evaluate in general but it is possible to use it in the special case where the spatial geometry is flat, exactly the situation we are discussing in this article.

Schoen and Yau define the size of any compact three dimensional subset (\( \Omega \)) of a Riemannian manifold as the minor radius of the largest three torus that can be embedded in \( \Omega \). They call this \( \text{Rad}(\Omega) \). In [19] they prove two theorems. Theorem I is a statement that one cannot have a large set with large positive scalar curvature. More precisely, they show that if the scalar curvature of \( \Omega \) is bounded below by a positive constant, \( (\sqrt{3} \mathcal{R})_0 > 0 \), then

\[ \text{Rad}(\Omega) \leq \sqrt{\frac{8 \pi^2}{3 \mathcal{R}_0}}. \]

Theorem II deals with the situation where one has a solution to the initial value constraints (including positive matter) on a set \( \Omega \). If the matter satisfies \( \rho - |J| \geq \lambda > 0 \) and if \( \Omega \) is large in the sense that

\[ \text{Rad}(\Omega) \geq \sqrt{\frac{3 \pi^2}{16 \pi \lambda}}, \]

then the initial data must have a trapped surface.

If we have a maximal slice, we have that \( (\sqrt{3} \mathcal{R}) = K_{ij} K_{ij} + 16 \pi \rho \geq 16 \pi \lambda \) and so we can use \( 16 \pi \lambda \) in place of \( \mathcal{R}_0 \) in Theorem I. However, since \( 8/3 < 3 \), we have from Theorem I that we can never get maximal initial data to satisfy Eq. (52). It is clear that the estimates leading to Theorem II are not sharp and some number smaller than 3 would almost certainly suffice. Unfortunately, the constant in Theorem I is also not sharp (see [20]) and the two constants are linked. Therefore our only hope of finding a nontrivial system in which to use Theorem II of Schoen and Yau is to look at nonmaximal initial data. In this case there is at least the possibility of having large \( \rho \), so as to satisfy the condition in Theorem II, while simultaneously having small \( (\sqrt{3} \mathcal{R}) \) to es-
cape the barrier that Theorem I imposes. When one looks at the way that Theorems I and II are derived, to try and find a configuration that satisfies the condition in Theorem II, it is clear that one wants to have no current, because it works against \( \rho \) and to have the mass density as uniform as possible. Further, one wants as little transverse traceless \((TT)\) part in the extrinsic curvature as possible so as to add to the scalar curvature. Finally, the metric should be as simple as possible. This leads one to consider the situation where the intrinsic metric is flat, and thus eliminating any barrier due to Theorem I; the extrinsic curvature is pure trace; and the trace is constant, so that there is no current and one has a constant mass density. The other great advantage of the flat metric is that one can easily evaluate their “torus” measure of the size of a set.

In our notation, this is equivalent to choosing \( K_R K = K \), or \( 0.5 \) and \( \text{Tr} K = \text{constant} \). From Eq. (2) this gives \( J = 0 \). From Eq. (7) we get that the mass density is constant and satisfies

\[
\rho = \rho_0 = \frac{3 K_R^2}{8 \pi}.
\]

If we have a spherical set of radius \( L \) satisfying this solution (one can think of it as part of a flat cosmology), it is clear from Eq. (30) that the horizon appears when \( |L K_R| = 1 \). The sufficient condition we have derived [Eq. (35)] when applied to this special case gives \( |L K_R| = \sqrt{2} \), and the necessary condition [Eq. (41)] gives \( |L K_R| > 0.84 \). Happily, these numbers lie on each side of 1.

If we apply the Schoen and Yau condition to an intrinsically flat constant density sphere of radius \( L \), we get that \( R \ell (\Omega) = L/2 \) and the Schoen and Yau sufficient condition [Eq. (52)] becomes \( |L K_R| = \sqrt{2} \). This calculation shows that the Schoen and Yau Theorem II is not vacuous. However, their set is 4 times larger than is required and their necessity condition a factor of 3 weaker than ours.

CONCLUSIONS

We have examined the constraints in spherically symmetric general relativity using an intrinsic time to foliate spacetime. Specifically we have foliated spacetime with flat spatial slices. The presentation of the initial data on a flat spatial hypersurface is very different from that on a hypersurface belonging to an extrinsic time foliation of spacetime. The Hamiltonian constraint becomes an algebraic constraint on the extrinsic curvature: when the weak energy condition is satisfied, trajectories in superspace lie completely inside the superspace light-cone—the complement in superspace of the allowed region consistent with any standard extrinsic curvature foliation \((K_R = 0\) or \(\text{Tr} K = 0\) or, more generally, any of those considered in [3]\).

In this foliation, solutions of the constraints do exhibit peculiarities: when the sources are finite, there are no singular geometries satisfying the constraints other than those which contain a singularity at their center; when the dominant energy condition is satisfied, \( K_R \) possesses a definite sign; they do not admit minimal surfaces. Despite this, we find that the physical description they provide of apparent horizons is completely consistent with that in an extrinsic time foliation. Not only do the natural measures of the material content for necessary and sufficient conditions \((\rho_{\text{Max}} \) and \( M \) respectively) coincide with those we found when we considered extrinsic time slices but, in addition, the inequalities assume identical forms.

Analogous gauges are applicable with other topologies. For example, in a closed cosmology with \( S^3 \) topology one could choose \( R(l) = (l_0/2\pi)\sin(\pi l/l_0) \), where \( l_0 \) is the inter-polar distance.

It would be interesting to examine the canonical reduction and subsequent quantization of spherically symmetric general relativity in this gauge. The fact that many of the features of extrinsic time foliations which are problematic do not occur suggests that flat foliations could provide a valuable alternative, in particular, for the description of gravitational collapse.

Finally, there appears to be no immediate obstruction to the construction of a foliation of a general asymptotically flat spacetime by a gauge of the form, \( 3 R = 0 \).

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[14] A further coordinate transformation to a freely falling coordinate system can be made by setting [6]

\[
R = \left( \frac{3}{2} \left( \frac{\rho}{\sqrt{2m}} + \sqrt{2m} \tau \right) \right)^{2/3}.
\]
Using $\xi$ as a spatial coordinate the Schwarzschild metric can be written as

$$
\begin{align*}
\text{ds}^2 &= -d\tau^2 + \left[ \frac{3}{2} \left( \frac{\xi}{\sqrt{2m}} + \sqrt{2m \tau} \right) \right]^{-2/3} d\Omega^2 \\
& \quad + \left[ \frac{3}{2} \left( \frac{\xi}{\sqrt{2m}} + \sqrt{2m \tau} \right) \right]^{4/3} d\Omega^2. 
\end{align*}
$$

[18] We have seen that when $J=0$ the quasilocal mass coincides with the bare mass, $m=M$. In an extrinsic time foliation of spacetime, the difference $m-M$ has been shown to be strictly negative and thus provides a good measure of the binding energy of the spherically symmetric system [1]. When $J \neq 0$ on a flat slice, however, $m-M$ does not possess a definite sign because, from Eq. (21), we get that $(m-M)^2 = 4 \pi \rho J^2 K_R$, which can be either positive or negative. We conclude that $m-M$ is not a satisfactory measure of binding energy in this gauge. A challenge is to define a quantity with a wider validity.