Figure 1: The FOSH operator shown acting on the connection coefficients. Dot indicates zero.
Figure 2: The FOSH operator shown acting on the Weyl components.

Figure 3: The FOSH operator in 2+1 gravity shown acting on the components $N_{ab}$. 
Hyperbolic equations for vacuum gravity using special orthonormal frames

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Abstract

By adopting Nester’s 4–dimensional special orthonormal frames, the tetrad equations for vacuum gravity are put into explicitly causal and symmetric hyperbolic form, independent of any time slicing or other gauge or coordinate specialization.
0.1. Introduction

We have previously given a well set and causal exterior differential system for vacuum gravity, generated by a closed set of differential forms describing the immersion of 4-dimensional spacetime into a flat 10 dimensional space [1]. The orthonormal frame bundle of the latter has a canonical basis of 10 (translation) 1-forms $\omega^\mu$ and 45 (rotation) 1-forms $\omega^\mu_{\nu}$, satisfying the structure equations of the Lie group $ISO(10)$. Dividing the range $\mu, \nu = 1, \ldots, 10$ into two ranges $i, j = 1, \ldots, 4$ and $A, B = 5, \ldots, 10$, these are

\begin{align*}
d\omega^i + \omega^i_j \wedge \omega^j + \omega^i_B \wedge \omega^B &= 0 \\
d\omega^A + \omega^A_j \wedge \omega^j + \omega^A_B \wedge \omega^B &= 0 \\
d\omega^i_k + \omega^i_j \wedge \omega^j_k + \omega^i_B \wedge \omega^B_k &= 0 \\
d\omega^i_A + \omega^i_j \wedge \omega^j_A + \omega^i_B \wedge \omega^A_B &= 0 \\
d\omega^A_C + \omega^A_j \wedge \omega^j_C + \omega^A_B \wedge \omega^B_C &= 0.
\end{align*}

The exterior differential system is generated by 6 (immersion) 1-forms $\omega^A$, their closure 2-forms $d\omega^A = -\omega^A_i \wedge \omega^i$, and 4 closed 3-forms ensuring Ricci-flatness, namely, $R_{ij} \wedge \omega_k \varepsilon^{ijkl}$, where $\frac{1}{2}R^i_j = d\omega^i_j + \omega^i_k \wedge \omega^k_j = -\omega^i_A \wedge \omega^A_j$ defines the Riemann 2-forms.

The significant result is the calculation of the Cartan characteristic integers

$s = \{s_0, s_1, s_2, s_3\} = \{6, 6, 10, 8\}$.

This shows the solutions to be 25 dimensional, regular (i.e., in principle constructed from a series of Cauchy-Kowalewsky integrations) and causal (i.e., $s_4 = 0$, so the solutions are determined from suitable data set on 3 dimensional immersed manifolds). The solutions are involutary with respect to $\omega^i, \omega^i_j$ and $\omega^A_j$: these forms remain independent when pulled back into a solution manifold and can be adopted as a basis there.

The six basis forms $\omega^i_j$ and 15 basis forms $\omega^A_B$, which occur of course in the structure relations, do not appear explicitly in the exterior differential system, showing the solutions to be bundles having 21 dimensional fibers over a four dimensional base. Evidently this expresses arbitrary $O(4)$ (for Lorentzian, $O(3, 1)$) rotations of the tetrad frame $\omega^i$ and $O(6)$ rotations of the immersion co-frame $\omega^A$ at each point of the base. On a four dimensional cross-section all the forms can be expanded on the $\omega^i$ basis, the $\omega^i_j$ being a metric connection.
In Section 2 we give a new immersion exterior differential system for vacuum gravity that incorporates the higher dimensional special orthonormal frame (HSOF) conditions proposed by Nester [2] as a generalization of the special orthonormal frame conditions (SOF) in three dimensional Riemannian geometry [3] [4] [5]. Calculation of the Cartan characteristic integers for this exterior differential system for 4 dimensional Riemannian geometry shows that it is also well set and causal, so that, as Nester conjectured, such frames can be imposed without impediment in extended regions of vacuum spacetime. Solutions are now 19 dimensional, the fibers expressing only arbitrary $O(6)$ rotations of the co-frames.

These exterior differential systems include all integrability conditions for the determination of a metric on the base, or on any 4 dimensional cross section. In terms of base space coordinates $x^\mu \ (\mu, \nu = 1, 2, 3, 4)$, invertible matrices of functions $^i\lambda_\mu(x)$ and $^i\lambda^\mu(x)$ exist such that $^i\lambda_\mu_j^\lambda = \delta^i_j, \ (i, j = 1, 2, 3, 4)$. We introduce the Minkowski metric, i.e., $\eta^{ij} = \text{diag}(1,1,1,-1)$, to raise and lower left, or Lorentz, indices. Then the metric is given by $g^{\mu\nu} = \eta^{kl} \lambda^\mu \lambda^\nu$. Inserting $\omega_i^\lambda = ^i\lambda_\mu dx^\mu$ in the structure equations and in the exterior differential system gives the partial differential equations for coordinate components of the tetrad field.

In Section 3 we reformulate the higher dimensional special orthonormal frame system in orthonormal tetrad components, using $\eta^{ij}$. We use the dyadic formalism of references [6] and [7], in which the unit 4−vector field $^4\lambda^\mu$ is given a special meaning: it traces a congruence of timelike world lines, a 3 parameter “fluid” of point observers, to which physical interpretations of the 24 components of the connection are attributed. The three spacelike unit vectors $^a\lambda^\mu \ (a = 1, 2, 3)$ at each point complete a local orthonormal frame for the observer there. We expand the six 1−forms $\omega_j^i$ on the $\omega^a, \omega^4$ basis $(a, b = 1, 2, 3)$, and we expand the six 2−forms $R_j^i$ on the $\omega^a \wedge \omega^b, \omega^a \wedge \omega^4$ basis, subject to the conditions for Ricci-flatness.

The 24 connection components are grouped into the following ensembles: $3 \times 3$ dyadics $K$ and $N$, and $3−$vectors $a$ and $\omega$. The dyadic $K$ has components $K_{ab}$, and can be resolved into

$$K_{ab} = S_{ab} - \Omega_c \varepsilon_{acb},$$

where $S_{ab}$ is the symmetric rate-of-strain 3−tensor of the observer fluid, and $\Omega_c$ is its axial vector of vorticity. Or, in dyadic notation

$$K = S - \Omega \times 1,$$
where $l$ is the unit dyadic. The dyadic $N$ is formed from the nine spacelike Ricci rotation coefficients of the $\omega^a$ basis. It has components $N_{ab}$ and can similarly be resolved into a symmetric part and an axial vector $n_c$

$$N_{ab} = N_{ab}^{sym} - n_c \varepsilon_{acb}.$$ 

Again we write in dyadic notation,

$$N = N^{sym} - n \times l.$$ 

The vector $\omega$ has components $\omega^a$ and is the time-dependent angular velocity of the triad seen by an observer moving along $\omega^4$, with respect to a Fermi-propagated frame. Since $\omega$ is a standard notation for angular velocity this usage should in context not be difficult to distinguish from the $1-$forms $\omega^a$ we have used up to this point. (When necessary, we can write the components of the angular velocity as $\varpi_a$.) The vector $a$ has components $a^a$ and is the acceleration of the point observers, i.e., their departure from geodesic motion. The vectors $a$ and $\omega$ are in principle determined operationally by the point observers, using spring balances and supported spinning particles in the local frames. The ten components of the Weyl tensor yield symmetric tracefree dyadics $A$ and $B$, the “electric” and “magnetic” tidal fields. The quantities $A$, $B$, $K$, $N$, $a$, and $\omega$ are defined in terms of the $\lambda_{\mu,\nu}$ in [6].

The dyadic formalism is completed by the use of inner and outer $3-$dimensional multiplication ($\cdot$ and $\times$), and by convective derivatives (also known as unit derivations) in the timelike and spacelike directions ($\dot{}$ and $D$). Use of spatial covariant differentiation $\nabla$, related to $D$ by the spatial connection $N$, often is more efficient than the use of $D$. Detailed expositions of this formalism, including the various $3-$vector and dyadic relations, the relations of $D$ and $\nabla$ (involving $N$), and the 52 general first order dyadic differential equations for vacuum $4-$geometry, are to be found in references [6] and [7].

There are some notational changes of which the reader should be aware. The dyadic we now (and in [7]) denote by $N$ was in [6] denoted by $N^*$. For the dyadic $N$ in [6] (which was symmetric) one should now understand

$$N - (Tr N) l + n \times l$$

and for the vector $L$ in [6] one should now write $n$.

Section 3 also gives the 12 new equations specializing to $4-$dimensional special orthonormal frames. Six of the additional equations involve $\dot{a}$ and $\dot{\omega}$, which had
not otherwise appeared in the tetrad equations. This is also the case with the six conditions for the “Lorentzian frame” specialization (based on time slicing) discussed by van Putten and Eardley [8]. When time evolution of all dependent variables is explicit the causal structure of the system can be seen as resulting from local wave propagation subject to constraints expressing integrability conditions. The integrability properties of these sets of dyadic equations are briefly discussed in Section 3 in terms of Cartan’s reduced characters and Cartan’s test.

In Section 4 we present the complete set of 64 dyadic equations resulting from substitution of dyadic components in the structure equations and the HSOF exterior differential system of Section 2. There are 34 equations involving the timelike derivatives of the 24 connection coefficients and the 10 Weyl components, together with an additional 30 transverse equations in which no time derivatives appear. By forming appropriate linear combinations, the final equations have been arranged to show that they have first order symmetric hyperbolic (FOSH) structure similar to that of recent formalisms involving preferred time slices, designed for application to numerical gravity [9] [10] [11] [12] [13].

A number of known and new “hyperbolic reductions” of the vacuum field equations, and the choices of gauge they allow, have been surveyed by Friedrich [14]. He uses both tetrad formalism, and ADM variables based on preferred time slicing. The rôle of the Bianchi identities is emphasized: the equations for $A$ (in [14], $E$) and $B$, which propagate as spin$−2$ massless fields (cf., e.g., [6]), are given in FOSH form. The FOSH equations obtained in the present paper necessarily include these. Nester’s conditions on the tetrad components of the connection however leave no further freedom in choice of gauge (or frame), and seem not to have been previously used. They result in FOSH structure with constant coefficients, and force all the dependent variables to evolve along null cones.

An analysis using Cartan’s test shows that in the present case the observer fluid necessarily has vorticity, i.e., $\Omega \neq 0$, so it is not 3-space orthogonal, and therefore Nester’s special orthonormal frames cannot be based on a preferred slicing. Nevertheless, the relative simplicity of these equations may be advantageous. No lapse or shift variables or higher order derivatives have been introduced. Of course, useful related coordinates can still be found. Two are suggested immediately by the closed 2-forms in the exterior differential system, and we note in Appendix 1 that the dyadic conditions for a harmonic timelike coordinate are again of first order symmetric hyperbolic (FOSH) form, so can simply be added to our results. Further adopting three spacelike coordinates co-moving with the
point observers then allows an explicit line element to be written, given a solution of the FOSH equations.

A similar application of special orthonormal frames can be made for 2 + 1 gravity, again leading to a constant coefficient FOSH system. This is outlined in Appendix 2.

0.2. Higher dimensional special orthonormal frames

Using eight new variables \( y_i \) and \( z_i \) \( (i = 1, 2, 3, 4) \), we prolong the previously given immersion exterior differential system for vacuum gravity, with two additional closed 2-forms, two additional 3-forms, and their closure 4-forms. The 3-forms and 4-forms essentially define the \( y_i \) and \( z_i \) in terms of the connection forms (the \( \omega^i_j \) pulled back into the solutions), and the 2-forms require them to satisfy Dirac type partial differential equations. When expanded in tetrad components in Section 3 it can be verified that this is precisely Nester’s prescription. The generators of this exterior differential system in 63 dimensions are

\[
\omega^A \\
\omega^i_A \wedge \omega^j \\
(dy_i - \omega^j_i y_j) \wedge \omega^j \\
(dz_i - \omega^j_i z_j) \wedge \omega^j \\
(\omega^i_A \wedge \omega^k \wedge \omega^l + \frac{2}{3} y_i \omega^j \wedge \omega^k \wedge \omega^l) \varepsilon^{ijkl} \\
\omega_{ij} \wedge \omega^j \wedge \omega^k - \frac{1}{3} z_i \omega^j \wedge \omega^k \wedge \omega^j \varepsilon^{ijkl} \\
\omega^A_j \wedge \omega^k_A \wedge \omega^l \varepsilon^{ijkl} \\
(-\omega^i_A \wedge \omega^s_A \wedge \omega^k \wedge \omega^l + 2\omega^i_A \wedge \omega^s_A \wedge \omega^k \wedge \omega^l + \frac{2}{3} dy_i \wedge \omega^j \wedge \omega^k \wedge \omega^l - 2y_i \omega^s_A \wedge \omega^k_A \wedge \omega^j \varepsilon^{ijkl} \\
-\omega_{is} \wedge \omega^j_A \wedge \omega^k \wedge \omega^l + 2\omega_{ij} \wedge \omega^s_A \wedge \omega^k \wedge \omega^j - (\frac{1}{3} dz_i \wedge \omega^j \wedge \omega^k \wedge \omega^l - z_i \omega^s_A \wedge \omega^k_A \wedge \omega^j) \varepsilon^{ijkl}. \quad (1)
\]

Monte Carlo calculations \( \Pi \) of the Cartan characteristic integers yield

\[ s = \{6, 8, 14, 16\} , \]

therefore the exterior differential system is well set and causal; solutions are 19 dimensional, fibered over 4 dimensions. The \( \omega^A_B \) do not appear in the exterior
differential system, so the fibers express co-frame rotation. But the $\omega_j^i$ are explicitly present so the frames are specialized and determined up to a simultaneous rotation at every point.

If the two 2—forms are exact, two further variables can be added, say $\zeta$ and $\eta$, together with 1—forms
\[
d\zeta - y_i\omega^i,
\]
\[
d\eta - z_i\omega^i.
\]
These should be useful for introducing intrinsic coordinates into the HSOF formulation.

0.3. The dyadic components of the connection forms $\omega_j^i$ and Riemann forms $R_j^i$

We expand the $\omega_j^i$ on $\omega^i$, defining the 24 3—dyadic and 3—vector components of $K$, $N$, $\omega$, and $a$, summarized in the Introduction and described in detail in references [6] and [7]
\[
\omega_{12} = -\omega_{21} = N_{13}\omega^1 + N_{23}\omega^2 + N_{33}\omega^3 + \omega^3\omega^4
\]
\[
\omega_{23} = -\omega_{32} = N_{21}\omega^2 + N_{31}\omega^3 + N_{11}\omega^1 + \omega^1\omega^4
\]
\[
\omega_{31} = -\omega_{13} = N_{32}\omega^3 + N_{12}\omega^1 + N_{22}\omega^2 + \omega^2\omega^4
\]
\[
\omega_{14} = -\omega_{41} = -K_{11}\omega^1 - K_{21}\omega^2 - K_{31}\omega^3 + a^1\omega^4
\]
\[
\omega_{24} = -\omega_{42} = -K_{22}\omega^2 - K_{32}\omega^3 - K_{12}\omega^1 + a^2\omega^4
\]
\[
\omega_{34} = -\omega_{43} = -K_{33}\omega^3 - K_{13}\omega^1 - K_{23}\omega^2 + a^3\omega^4.
\]
We also expand the Riemann 2—forms on a basis consisting of $\omega^a \wedge \omega^4$ and $\omega^a \wedge \omega^b (a, b = 1, 2, 3)$, to define (in the Ricci-flat case) Weyl dyadics $A$ and $B$
\[
\frac{1}{2}R_{21} = B_{31}\omega^1 \wedge \omega^4 + B_{32}\omega^2 \wedge \omega^4 + B_{33}\omega^3 \wedge \omega^4 + A_{31}\omega^2 \wedge \omega^3 + A_{32}\omega^3 \wedge \omega^1 + A_{33}\omega^1 \wedge \omega^2
\]
\[
\frac{1}{2}R_{32} = B_{12}\omega^2 \wedge \omega^4 + B_{13}\omega^3 \wedge \omega^4 + B_{11}\omega^1 \wedge \omega^4 + A_{12}\omega^3 \wedge \omega^1 + A_{13}\omega^1 \wedge \omega^2 + A_{11}\omega^2 \wedge \omega^3
\]
\[
\frac{1}{2}R_{13} = B_{23}\omega^3 \wedge \omega^4 + B_{21}\omega^1 \wedge \omega^4 + B_{22}\omega^2 \wedge \omega^4 + A_{23}\omega^1 \wedge \omega^2 + A_{21}\omega^2 \wedge \omega^3 + A_{22}\omega^3 \wedge \omega^1
\]
\[
\frac{1}{2}R_{41} = -A_{11}\omega^1 \wedge \omega^4 - A_{12}\omega^2 \wedge \omega^4 - A_{13}\omega^3 \wedge \omega^4 + B_{11}\omega^2 \wedge \omega^3 + B_{12}\omega^3 \wedge \omega^1 + B_{13}\omega^1 \wedge \omega^2
\]
\[
\frac{1}{2}R_{42} = -A_{22}\omega^2 \wedge \omega^4 - A_{23}\omega^3 \wedge \omega^4 - A_{21}\omega^1 \wedge \omega^4 + B_{22}\omega^3 \wedge \omega^1 + B_{23}\omega^1 \wedge \omega^2 + B_{21}\omega^2 \wedge \omega^3
\]
\[
\frac{1}{2}R_{43} = -A_{33}\omega^3 \wedge \omega^4 - A_{31}\omega^1 \wedge \omega^4 - A_{32}\omega^2 \wedge \omega^4 + B_{33}\omega^1 \wedge \omega^2 + B_{31}\omega^2 \wedge \omega^3 + B_{32}\omega^3 \wedge \omega^1.
\]
The Riemann 2-forms given above are antisymmetric under interchange of their indices. The 10 Weyl components satisfy $A_{ab} = A_{ba}$, $A_{aa} = 0$, $B_{ab} = B_{ba}$, and $B_{aa} = 0$. It can be verified that the Riemann symmetries are satisfied, namely,

$$R^a_b \wedge \omega^b + R^a_d \wedge \omega^d = 0$$

and

$$R^d_d \wedge \omega^b = 0,$$

and also that the four 3-form conditions for Ricci-flatness, namely,

$$R^i_j \wedge \omega^k \varepsilon^j_{ikl} = 0$$

have been imposed.

From the 3-forms of (1) we can now find the component expansions of the fields $y_i$ and $z_i$, on an integral manifold of the exterior differential system. These are given by the following

$$y_i = (y^a, -y^4) = (N \dot{x}l - a, +Tr K)$$

$$z_i = (z^a, -z^4) = (-K \dot{x}l + \omega, Tr N).$$

We could alternatively have used vectors $2n = N \dot{x}l$ and $2\Omega = K \dot{x}l$. The twelve first order dyadic equations arising from the two additional HSOF 2-forms then are

$$+(N \dot{x}l - a)^\cdot \nabla (Tr K) + (Tr K) a + K \cdot (N \dot{x}l - a) + \omega \times (N \dot{x}l - a) = 0$$

$$\nabla \times (N \dot{x}l - a) + (Tr K) K \dot{x}l = 0$$

$$+(K \dot{x}l - \omega)^\cdot \nabla (Tr N) - (Tr N) a + K \cdot (K \dot{x}l - \omega) + \omega \times (K \dot{x}l) = 0$$

$$\nabla \times (K \dot{x}l - \omega) - (Tr N) K \dot{x}l = 0.$$
\(\omega^i, \omega^i \wedge \omega^j, \text{etc.}\). So-called reduced characters \(s'_i\) are then conveniently computed from the left-hand sides alone, and Cartan’s test is to calculate

\[
h = \sum_{i=0}^{g-1} (g - i) \cdot s'_i.
\]

If \(h\) is equal to the number of independent first order equations, this establishes involutivity, i.e., the well set nature of the problem. Moreover, if

\[
s'_g = n - g - \sum_{i=0}^{g-1} s'_i = 0,
\]

the Cauchy-Kowalewsky solutions are causal, i.e., determined by data given on a \(g - 1\) dimensional surface. In the present context \(g = 4\).

The 52 general vacuum dyadic equations were of this form \[6\] \[15\]. They result from computing the exterior derivatives of (2) and (3), so have six 2–forms and six 3–forms in 34 dependent variables. Their left hand sides are

\[
\begin{align*}
&dN_{ba} \wedge \omega^b - d\mathcal{W}_a \wedge \omega^4 \\
&dK_{ba} \wedge \omega^b + da_a \wedge \omega^4 \\
&dA_{ab} \wedge \omega^b \wedge \omega^4 - \frac{1}{6} dB_{ab} \wedge \omega^c \wedge \omega^d \varepsilon^{bc}_{\cdot cd} \\
&dB_{ab} \wedge \omega^b \wedge \omega^4 + \frac{1}{6} dA_{ab} \wedge \omega^c \wedge \omega^d \varepsilon^{bc}_{\cdot cd}.
\end{align*}
\]

We calculate \(s' = \{0, 6, 12, 10\}\), so \(h = 52\). There are however \(s'_4 = 6\) arbitrary functions in the solution, and this system is therefore not causal. The new HSOF system of dyadic equations adjoins 12 equations in two additional 2–forms to express equations (6)–(9). Their left hand sides are

\[
\begin{align*}
&d(2n_a - a_a) \wedge \omega^a + d(Tr K) \wedge \omega^4 \\
&d(2\Omega_a - \mathcal{W}_a) \wedge \omega^a - d(Tr N) \wedge \omega^4.
\end{align*}
\]

The reduced characters are \(s' = \{0, 8, 14, 12\}\), \(h = 64\), \(s'_4 = 0\), and therefore the final system is causal.
0.4. The FOSH dyadic equations

Linear combinations of the 52 general dyadic equations and the 12 new conditions due to Nester can now easily be made to put the result in FOSH form. The result is 34 equations involving the time derivatives and symmetric space derivatives of the dyadic components and $A$ and $B$, and 30 “constraint” or transverse relations not involving time derivatives. We give them in the following, written in full with their quadratic right hand sides. For the connection components we obtain

$$\dot{a} - \nabla \cdot K^T + \nabla \times \omega = -K \times N - \omega \cdot N + (Tr N) \omega + 2K \cdot n \quad (10)$$

$$\dot{\omega} + \nabla \cdot N^T - \nabla \times a = -(Tr N) a + (2\Omega - \omega) \cdot K - 2(Tr K) \Omega - 2n \cdot N \quad (11)$$

$$\dot{K} - a \nabla - 2\nabla n + 2n \nabla = -K \cdot K - 2(Tr K) \Omega \times I - \omega \times K + K \times \omega + aa - A \quad (12)$$

$$\dot{N} + \omega \nabla + 2\nabla \Omega - 2\Omega \nabla = -K \cdot N - \omega \times N + K \times a - a\omega - 2(Tr N) \Omega \times I + B. \quad (13)$$

We mention in passing that using $D$ instead of $\nabla$ usually makes the right hand sides of these equations less concise. The prominent exception to this statement is the first equation, which becomes homogeneous and linear, namely,

$$\dot{a} - D \cdot K^T + D \times \omega = 0.$$  

It has been convenient to use $n$ and $\Omega$, and $\nabla$ can operate from the right as well as from the left, to express transposed indices. It is best to return to the nine components of $K$ and of $N$ to see the FOSH structure; e.g. write $(2\nabla n)_{11} = \nabla_1 N_{23} - \nabla_1 N_{32}, \ (2n \nabla)_{23} = \nabla_3 N_{31} - \nabla_3 N_{13}, \ (2\Omega \nabla)_{22} = \nabla_2 K_{31} - \nabla_2 K_{13},$ etc. The left hand sides of the FOSH equations (10)–(13) are shown in Figure 1 as the product of a linear operator and a column vector.

The 24 constraint equations are

$$\nabla \times K = 2\Omega a - B \quad (14)$$

$$\nabla \times N = -\frac{1}{2} N^T \times N - \frac{1}{2} K \times K + (\Omega \cdot K) \times I - 2\Omega \omega - A \quad (15)$$

$$\nabla \times (a - 2n) = 2(Tr K) \Omega \quad (16)$$

$$\nabla \times (\omega - 2\Omega) = -2(Tr N) \Omega \quad (17)$$

By taking a second time derivative of the FOSH equations, interchanging space and time derivatives, and substituting back both the FOSH and the constraint
equations, it can be seen that the dyadic variables all propagate causally along null cones.

The 10 FOSH equations for the Weyl components are those for traceless transverse massless spin$−2$ fields (dyadic and Bianchi identities, linearly combined) and are as follows

$$2 \dot{B} - \nabla \times A + A \times \nabla =
2B \times \omega - 2\omega \times B - 2A \times a + 2a \times A + K^\tau \cdot B + B \cdot K - 2(Tr K)B + K \times B + B \times K \tag{18}$$

$$2 \dot{A} + \nabla \times B - B \times \nabla =
2A \times \omega - 2\omega \times A + 2B \times a - 2a \times B + K^\tau \cdot A + A \cdot K - 2(Tr K)A + K \times A + A \times K \tag{19}$$

The left hand sides of the Bianchi equations are shown in Figure 2 as the product of a linear operator and a column vector.

The final 6 constraint equations are

$$\nabla \cdot A = -K \times B - 4\Omega \cdot B \tag{20}$$

$$\nabla \cdot B = K \times A + 4\Omega \cdot A. \tag{21}$$

Finally, it should be remarked that our derivations from well set exterior differential systems obviate any need to verify that the transverse (or constraint) equations, i.e., those not involving time derivatives, are compatible with the FOSH system, and that they are propagated invariantly by it.

0.5. Appendix 1: Harmonic and co-moving coordinates and an explicit line element

The formulation of this paper is coordinate free and gauge independent. It may however also be of use to briefly record how a harmonic time coordinate, that is, one which satisfies a wave equation, and co-moving (with $4\lambda^\mu$) spacelike coordinates can be adopted.

In coordinate language put $t_{\mu\nu} g^{\mu\nu} = 0$ ($\mu, \nu = 1, \cdots, 4$). Introducing fields $\phi = \dot{t}$, $A = \nabla t$, the first order tetrad equations for this are

$$\dot{A} - \nabla \phi = -K \cdot A - \omega \times A + \phi a$$
\[ \dot{\phi} - \nabla \cdot A = a \cdot A - \dot{\phi} (Tr K) \]
\[ \nabla \times A = 2 \phi \Omega. \]

This set consists of four FOSH equations plus three constraints. It can be added to, and solved simultaneously with, the equations in Section 4.

Co-moving spacelike coordinates, say \( \dot{x}^\alpha (\alpha = 1, \cdots, 3) \), such that \( \dot{x}^\alpha = 0 \) are found by setting
\[ e^\alpha = \nabla x^\alpha, \]
and the integrability conditions for this are
\[ \dot{e}^\alpha = -K \cdot e^\alpha - \omega \times e^\alpha \]
\[ \nabla \times e^\alpha = 0. \]

The \( e^\alpha \) are, understandably, the first dependent variables we have found whose causal propagation is strictly timelike and not along the local null cone. We introduce coordinate components \( A \cdot e^\alpha = A^\alpha \) and \( h^{\alpha \beta} = e^\alpha \cdot e^\beta \), and also calculate the inverse \( h_{\alpha \beta} \) (\( h^{\alpha \gamma} h_{\alpha \beta} = \delta_\beta^\gamma \)). Then \( \phi, A_\alpha = h_{\alpha \beta} A^\beta \) and \( h_{\alpha \beta} \), functions of \( x^\alpha \), and \( t \), enter the final line element:
\[ ds^2 = -\phi^{-2} dt^2 + 2\phi^{-2} A_\alpha dx^\alpha dt + (h_{\alpha \beta} - \phi^{-2} A_\alpha A_\beta) dx^\alpha dx^\beta. \]

In covariant 4—vector terms, the coordinate components of the tetrad vectors take the form
\[ r^\lambda_\mu = \begin{pmatrix} a^\lambda A & a \lambda A \\ 0 & \phi \end{pmatrix} \]
\[ r^\lambda_\mu = \begin{pmatrix} a^\lambda A & 0 \\ -\phi^{-1} A_\lambda & \phi^{-1} \end{pmatrix} \]

The orthonormal triad components of the vector \( A \) are written \( {}^a_\alpha A_\alpha \).

### 0.6. Appendix 2: 2 + 1 dimensional gravity

An entirely parallel development can be made using special orthonormal frames (SOF) in 2+1 gravity. The exterior differential system for immersion of 3—dimensional flat spaces in the 21 dimensional orthonormal frame bundle over 6 dimensional flat space, which is generated by the usual immersion forms \( \omega^A \) and \( d\omega^A (A = 4, 5, 6) \),
and the Riemann 2–forms $R^a_{bc} (a, b, c = 1, 2, 3)$, has $s = \{3, 6, 3\}$ and $g = 9$. We add a 2–form

$$
\omega^b_a \wedge \omega^c_{bc} + y_a \omega^b \wedge \omega^c_{bc},
$$

a 3–form

$$
\omega_{ab} \wedge \omega^a \wedge \omega^b - z \omega^1 \wedge \omega^2 \wedge \omega^3,
$$

and their closures, to introduce 4 new variables $y_a$ and $z$, so the exterior differential system is prolonged to a total of 25 dimensions. The SOF conditions are imposed as two additional closed exterior forms, namely,

$$
dz
$$

$$
d(y_a \omega^a).
$$

These adjoin solutions of a linear Dirac equation [4]. Now the Cartan characters show the exterior differential system to be well set and causal with $s = \{4, 8, 7\}$, and $g = 6$. Solutions are 3–parameter co-frame rotation bundles over 3–space, determined by causal integration from 2–spaces.

Orthonormal triad formalism for 3–dimensional geometry yields 9 equations for 9 Ricci rotation coefficients grouped in a dyadic $N$:

$$
\nabla \times N + \frac{1}{2} N^T \times N = 0.
$$

The SOF conditions require the identification of $z$ as $Tr N$ and the $y_a$ as the components of $n = \frac{1}{2} N \times 1$. To the nine equations for a flat 3–space are added six SOF equations:

$$
\nabla (Tr N) = 0
$$

$$
\nabla \times n = 0.
$$

Inserting a $-1$ corresponding to timelike nature of the 3 direction, the complete set falls into FOSH form with 6 transverse equations. The left hand sides of the FOSH equations are shown in Figure 3 as the product of a linear operator and a column vector.

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Bibliography

[1] Estabrook F B and Wahlquist H D 1993 Immersion ideals and the causal structure of Ricci-flat geometries Class. Quantum Grav. 10 1851-8

[2] Nester J M 1992 Special orthonormal frames J. Math. Phys. 33 910-13

[3] Nester J M 1991 Special orthonormal frames and energy localization Class. Quantum Grav. 8 L 19-23

[4] Nester J M 1989 A gauge condition for orthonormal three-frames J. Math. Phys. 30 624-6

[5] Dimakis A and Müller-Hoissen F 1989 On a gauge condition for orthonormal 3−frames Phys. Lett. A 142 73-4

[6] Derivation of the dyadic formalism, explicit expression in covariant coordinates, and physical interpretations in terms of local Lorentz observers can be found in Estabrook F B and Wahlquist H D 1964 Dyadic analysis of spacetime congruences J. Math. Phys. 5 1629-44

[7] Wahlquist H D 1992 The problem of exact interior solutions for rotating rigid bodies in general relativity J. Math. Phys. 33 304-35. The dyadic notation and general dyadic equations are given in appendices A and B. Errata for the paper as a whole are in J. Math. Phys. 33 3255, although appendices A and B are unchanged.

[8] van Putten M H P M and Eardley D M 1996 Nonlinear wave equations for relativity Phys. Rev. D 53 3056-63

[9] Choquet-Bruhat Y and York J W Jr 1995 Geometrical well-posed system for the Einstein equations C. R. Acad. Sci. Paris 321 Serie I 1089-95
[10] Abrahams A, Anderson A, Choquet-Bruhat Y and York J W Jr 1995 Einstein and Yang Mills in hyperbolic form without gauge fixing Phys. Rev. Lett. 75 3377-81

[11] Abrahams A, Anderson A, Choquet-Bruhat Y and York J W Jr 1997 Geometrical hyperbolic systems for general relativity and gauge theories Class. Quantum Gravity 14 A 9-22

[12] Bona C and Massó J 1989 Einstein evolution equations as a system of balance laws Phys. Rev. D 40 1022

[13] Bona C and Massó J 1992 Hyperbolic evolution system for numerical relativity Phys. Rev. Lett. 68 1097

[14] Friedrich H 1996 Hyperbolic reductions for Einstein’s equations Class. Quantum Gravity 13 1451-69. Also cf. Frittelli S and Reula O A 1996 First order symmetrical hyperbolic Einstein equations with arbitrary fixed gauge Phys. Rev. Lett. 76 4667-70

[15] Estabrook F B and Unti T W J 1968 Exterior forms and general relativity JPL Space Programs Summary 37-53, III, 158-60