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A NEW TENSORIAL CONSERVATION LAW FOR MAXWELL FIELDS ON THE KERR BACKGROUND

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Abstract. A new, conserved, symmetric tensor field for a source-free Maxwell test field on a four-dimensional spacetime with a conformal Killing-Yano tensor, satisfying a certain compatibility condition, is introduced. In particular, this construction works for the Kerr spacetime.

1. Introduction

In this paper, we consider the Maxwell equation for a real 2-form \( F_{ab} = F_{[ab]} \),
\[
\nabla^a F_{ab} = 0, \quad \nabla^a * F_{ab} = 0, \tag{1.1}
\]
on a four-dimensional Lorentzian manifold \((M, g_{ab})\). Recall that a conformal Killing-Yano tensor is a 2-form \( Y_{ab} = Y_{[ab]} \) satisfying
\[
\nabla^a (Y_b^c) = - \frac{1}{3} g_{ab} \nabla^d Y_{cd} + \frac{1}{3} g_{(a|c}] \nabla^d Y_{b)d}. \tag{1.2}
\]
Associated with \( Y_{ab} \) is the complex 1-form
\[
\xi_a = \frac{i}{3} \nabla_b Y_{ab}^c - \frac{i}{3} \nabla_b Y_{ac}^c. \tag{1.3}
\]
We say that \( Y_{ab} \) satisfies the aligned matter condition if the Ricci curvature and \( Y_{ab} \) satisfy
\[
R_{(a} Y_{b)c} = 0, \quad R_{(a} Y_{b)c}^* = 0. \tag{1.4}
\]

Theorem 1.1. Let \( Y_{ab} \) and \( F_{ab} \) be real 2-forms. Define the real 2-form \( Z_{ab} \) and the complex 1-form \( \eta_a \) by
\[
Z_{ab} = - \frac{4}{3} (*) F_{[a} Y_{b)c], \tag{1.5}
\]
\[
\eta_a = - \frac{1}{3} \nabla_b Z_{ab}^c - \frac{1}{3} i \nabla_b Z_{ab}^c, \tag{1.6}
\]
and the real symmetric 2-tensor \( V_{ab} \) by
\[
V_{ab} = \eta_{(a} \overline{\eta}_{b)} - \frac{1}{2} g_{ab} \xi - \frac{1}{3} (\mathcal{L}_{\text{Re} \xi} F)_{(a} \overline{Z}_{b)c} + \frac{1}{12} g_{ab} (\mathcal{L}_{\text{Re} \xi} F)^{cd} Z_{cd}
+ \frac{1}{3} (\mathcal{L}_{\text{Im} \xi} F)_{(a} Z_{b)c} - \frac{1}{12} g_{ab} (\mathcal{L}_{\text{Im} \xi} F)^{cd} Z_{cd}, \tag{1.7}
\]
where \( \xi_a \) is given by equation (1.3) and \( \overline{\eta}_a \) denotes the complex conjugate of \( \eta_a \).

If \( Y_{ab} \) is a conformal Killing-Yano tensor satisfying the aligned matter condition (1.4) and \( F_{ab} \) satisfies the Maxwell equations (1.1), then \( V_{ab} \) has vanishing divergence, \( \nabla^a V_{ab} = 0 \).

Remark 1.2. 1. The vector field \( \xi^a \) is Killing, \( \nabla_{(a} \xi_{b)} = 0 \), if the aligned matter condition (1.4) holds, cf. equation (2.9) below. If \( \nabla^a Y_{ab} = 0 \) then \( Y_{ab} \) is a Killing-Yano tensor, in which case \( \xi_a \) is real, and the last two terms of (1.7) vanish.

2. The Kerr family of stationary, rotating vacuum black hole metrics admit a Killing-Yano tensor. More generally, the Kerr-Newman family of stationary, rotating electro-vacuum black hole metrics admit a Killing-Yano tensor satisfying the aligned matter condition. See section 3 for further discussion.

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Let
\[ T_{ab} = -F_{ac}^\epsilon F_{bc}^\epsilon + \frac{1}{4} g_{ab} F_{cd} F^{cd} \]
be the symmetric energy-momentum tensor for the Maxwell field. It is traceless and satisfies the dominant energy condition, i.e. \( T_{ab} \mu^a \nu^b \geq 0 \) for any future causal vectors \( \mu^a, \nu^b \). Further, if \( F_{ab} \) satisfies the Maxwell equations, \( T_{ab} \) is conserved, \( \nabla^a T_{ab} = 0 \). Hence, the current
\[ J_a = T_{ab} \nu^b \]
(1.8)
is conserved, \( \nabla^a J_a = 0 \), if \( \nu^a \) is a conformal Killing field, \( \nabla^a (\nu_b - \frac{1}{4} g_{ab} \nu^c g_{bc} = 0 \).

For the Maxwell field on Minkowski space, and more generally on spacetimes admitting conformal Killing-Yano tensors satisfying the aligned matter condition, there are non-classical conserved currents not equivalent to any of the classical conserved energy-momentum currents of the form (1.8), see [1] and references therein. For the Maxwell field on Minkowski space, these include chiral currents constructed using the 20-dimensional family of conformal Killing-Yano tensors of Minkowski space. As shown by the authors [2], analogous conserved currents exist also on spacetimes with conformal Killing-Yano tensors satisfying the aligned matter condition.

In spite of the large literature on conformal Killing-Yano tensors, and the related conservation laws, the tensorial conservation law exhibited in Theorem 1.1 appears to be new, even in the Minkowski case. The fact that the new higher order tensor concomitant \( V_{ab} \) is conserved also in the case of the Kerr and Kerr-Newman spacetimes makes it interesting from the point of view of the black hole stability problem, which in fact served as an important motivation for the investigation which led to its discovery. See section 3 below for further remarks.

At this point, we should mention that the symmetric tensor
\[ B_{ab} = \nabla^d F_{bc} \nabla^d F_{ac} - \frac{1}{4} g_{ab} \nabla^f F_{cd} \nabla^f F^{cd} \]
which arises as a trace of the 4-index Chevreton tensor, was shown by Bergqvist et al. [6] to be traceless and conserved for a Maxwell field on a Ricci flat spacetime. Like the conserved tensor \( V_{ab} \) introduced in this paper, the tensor \( B_{ab} \) introduced by Bergqvist et al. depends on the \( F_{ab} \) and its first derivatives. However, while \( B_{ab} \) is traceless and fails to satisfy any positivity condition, the new tensor \( V_{ab} \) has trace \( V^a_a = -\eta^a \bar{\eta}_a \) and satisfies a weak form of the dominant energy condition in the sense the its leading order term, \( \eta_a \bar{\eta}_b = \frac{1}{2} g_{ab} \eta^c \bar{\eta}_c \) which is quadratic in first derivatives of \( F_{ab} \), is a superenergy tensor for \( \eta_a \) and hence does satisfy the dominant energy condition. Hence, energies can be constructed in terms of \( V_{ab} \) which are non-negative up to terms of lower order.

The proof of Theorem 1.1, which will be given in the next section, makes use of computations in the 2-spinor formalism. In the investigations leading to the main result, the SymManipulator package [3], developed by one of the authors (T.B.) for the Mathematica based symbolic differential geometry suite xAct [4], has played an essential role. SymManipulator makes it possible to systematically exploit decompositions in terms of irreducible representations of the spin group \( \text{SL}(2, \mathbb{C}) \), and allows one to carry out investigations that are not feasible by hand.

In section 3 we show how the main result applies for the Kerr-Newman family of electro-vacuum spacetimes, and indicate its relation to the Teukolsky and Teukolsky-Starobinsky equations.

---

1 A conserved current \( J_a \) is a 1-form concomitant of the Maxwell field, satisfying \( \nabla^a J_a = 0 \). We say that \( J_a \) is equivalent to \( \tilde{J}_a \) if \( J_a - \tilde{J}_a = \nabla^b C_{ab} \) for some 2-form \( C_{ab} = C_{(ab)} \).
2. Proof of theorem 1.1

For the remainder of this paper, we will make use of the 2-spinor formalism, following the conventions of [8]. Since our considerations are local, we can assume without loss of generality that \((M, g_{ab})\) is oriented and globally hyperbolic. This also implies that \(M\) is spin.

The spin group is SL(2, C) which has the inequivalent spinor representations \(C^2\) and \(\mathcal{C}^2\). Unprimed upper case latin indices and their primed versions are used for sections of the corresponding spinor bundles, respectively. The correspondence between spinors and tensors makes it possible to translate all tensor expressions to spinor form. The action of SL(2, C) on \(C^2\) leaves invariant the spin metric \(\epsilon_{AB} = \epsilon_{[AB]}\), which is used to raise and lower indices on tensors. The metric \(g_{ab}\) is related to \(\epsilon_{AB}\) by \(g_{ab} = \epsilon_{AB} \epsilon^{AB}\). Let \(S_{k,l}\) denote the space of symmetric spinors with \(k\) unprimed indices and \(l\) primed indices.

There are symmetric spinors \(\kappa_{AB}, \phi_{AB}\), and \(\Theta_{AB}\) such that

\[
Y_{ab} = \frac{1}{2i}(\tilde{\epsilon}_{AB}^{C} \kappa_{AB} - \epsilon_{AB} \kappa_{A'B'}),
\]

\[
F_{ab} = \epsilon_{AB} \phi_{AB} + \epsilon_{AB} \phi_{A'B'},
\]

\[
Z_{ab} = \epsilon_{AB} \Theta_{AB} + \epsilon_{AB} \bar{\Theta}_{A'B'}.
\]

The normalization of \(Y_{ab}\) is chosen for convenience. Equations (1.1)-(1.7) become respectively

\[
\nabla A_{AB} = 0,
\]

(2.1)

\[
\nabla_{(A} \kappa_{AB)BC} = 0,
\]

(2.2)

\[
\xi_{A} A'B = \nabla B_{A} \kappa_{AB},
\]

(2.3)

\[
\Phi_{(A} C_{|A'B'|B)C} = 0.
\]

(2.4)

\[
\Theta_{AB} = -2 \kappa_{A} C_{|B|C},
\]

(2.5)

\[
\eta_{A} A' = \nabla B_{A} \Theta_{AB}.
\]

(2.6)

and

\[
V_{AB} A'B' = \frac{1}{2} \eta_{AB} \tilde{\eta}_{A'B'} + \frac{1}{2} \eta_{BA} \eta_{B'A'} + \frac{1}{4} \Theta_{AB} (\tilde{\Theta}_{\tilde{A}} \tilde{\tilde{\Theta}}_{\tilde{B}}) A'B' + \frac{1}{4} \Theta_{A'B'} (\tilde{\Theta}_{\tilde{A}} \tilde{\tilde{\Theta}}_{\tilde{B}})_{AB},
\]

(2.7)

where \(\tilde{\Theta}_{\tilde{A}}\) is a conformally weighted Lie derivative on spinors, see equation (2.10) below.

The projection of the spinor covariant derivative \(\nabla_{AA'}\) on symmetric spinors (which form the irreducible representations of the spin group SL(2, C)) gives the following fundamental operators.

Definition 2.1 ([H Definition 13]). Let the differential operators \(\mathcal{D}_{k,l} : S_{k,l} \rightarrow S_{k-1,l-1}, \mathcal{C}_{k,l} : S_{k,l} \rightarrow S_{k+1,l-1}, \mathcal{G}_{k,l} : S_{k,l} \rightarrow S_{k-1,l+1}, \text{ and } \mathcal{F}_{k,l} : S_{k,l} \rightarrow S_{k+1,l+1}\) be defined by

\[
(\mathcal{D}_{k,l} \varphi)_{A_{1} \ldots A_{k-1} A'_{1} \ldots A'_{l-1}} = \nabla B_{A'_{1} \ldots A'_{l-1}} B'_{A_{1} \ldots A_{k-1}},
\]

(3.1)

\[
(\mathcal{C}_{k,l} \varphi)_{A_{1} \ldots A_{k+1} A'_{1} \ldots A'_{l-1}} = \nabla (A'_{1} B'_{A1} \ldots A_{k+1}) B'_{A_{1} \ldots A_{k-2}},
\]

(3.2)

\[
(\mathcal{G}_{k,l} \varphi)_{A_{1} \ldots A_{k-1} A'_{1} \ldots A'_{l+1}} = \nabla (A'_{1} A_{1} \ldots A_{k-1} B_{A_{1} \ldots A_{k+1}}),
\]

(3.3)

\[
(\mathcal{F}_{k,l} \varphi)_{A_{1} \ldots A_{k+1} A'_{1} \ldots A'_{l+1}} = \nabla (A'_{1} A_{1} \ldots A_{k-1} B_{A_{1} \ldots A_{k+1}}),
\]

(3.4)

The operators are called respectively the divergence, curl, curl-dagger, and twistor operators.

With respect to complex conjugation, the operators \(\mathcal{D}, \mathcal{F}\) satisfy \(\mathcal{D}_{k,l} = \mathcal{D}_{l,k}, \mathcal{F}_{k,l} = \mathcal{F}_{l,k}\), while \(\mathcal{G}_{k,l} = \mathcal{G}_{l,k}^{\dagger}, \mathcal{C}_{k,l} = \mathcal{C}_{l,k}^{\dagger}\). In the following, we shall use the
fundamental operators and their properties freely. Any covariant expression in spinors and their covariant derivatives can be written in terms of the fundamental operators using the following Lemma.

**Lemma 2.2** ([4] Lemma 15). For any \( \varphi_{A_1 \ldots A_k} A'_1 \ldots A'_{l+1} \in S_{k,l} \), we have the irreducible decomposition

\[
\nabla_{A_1} A'_1 \varphi_{A_2 \ldots A_{k+1}} A'_2 \ldots A'_{l+1} + \text{terms} = (\mathcal{T}_{k,l}\varphi)_{A_1 \ldots A_{k+1}} A'_1 \ldots A'_{l+1} + \text{terms},
\]

where the \( \varphi \) is related to the fundamental operators by the fundamental operators, see [4, Lemma 18]. The following lemma gives the decomposition aligned matter condition (2.4)

\[L \cdot \text{operators using the following Lemma.}\]

Directly from the Killing spinor equation and the commutators (2.8a) and (2.8d)

\[
\text{we get}
\]

\[
(\mathcal{P}_{AB})_{AB} = 0,
\]

and

\[
(\mathcal{G}_{2,0})_{ABCA'} = 0
\]

respectively, in terms of the fundamental operators.

In the computations below we shall need some commutator relations satisfied by the fundamental operators, see [4] Lemma 18. The following lemma gives the commutators which are relevant here.

**Lemma 2.3.** Let \( \varphi_{AB} \in S_{2,0} \). The operators \( \mathcal{D}, \mathcal{G}, \mathcal{C}^\dagger \) and \( \mathcal{T} \) satisfies the following commutator relations

\[
(\mathcal{D}_{1,1} \mathcal{C}_{2,0}^\dagger \varphi) = 0,
\]

\[
(\mathcal{G}_{3,1} \mathcal{T}_{2,0} \varphi)_{ABCD} = 2\Psi_{(ABC}^F \varphi_{DF)}
\]

\[
(\mathcal{G}_{3,1} \mathcal{T}_{2,0} \varphi)_{AB'A'B'} = \frac{2}{3} (\mathcal{P}_{1,1} \mathcal{C}_{2,0}^\dagger \varphi)_{AB'A'B'} + 2\Phi_{(A}^C |A'B'| \varphi_{BC)}
\]

\[
(\mathcal{P}_{3,1} \mathcal{T}_{2,0} \varphi)_{AB} = - \frac{2}{3} (\mathcal{P}_{1,1} \mathcal{C}_{2,0}^\dagger \varphi)_{AB} - 8\Lambda \varphi_{AB} + 2\Psi_{ABCD} \varphi_{CD}.
\]

Directly from the Killing spinor equation and the commutators (2.8a) and (2.8d)

we get

\[
(\mathcal{P}_{1,1} \xi) = 0,
\]

\[
(\mathcal{P}_{1,1} \xi)_{AB} = - 3\Phi_{(A}^C |A'B'| \xi_{BC)}.
\]

Hence, if the aligned matter condition is satisfied, \( \xi^{AA'} \) is a Killing vector.

Given a conformal Killing vector \( \xi^{AA'} \), we define a conformally weighted Lie derivative acting on a symmetric valence \((2s,0)\) spinor field by [4] Definition 17

\[
\mathcal{L}_\xi \varphi_{A_1 \ldots A_{2s}} = \xi^{BB'} \nabla_{BB'} \varphi_{A_1 \ldots A_{2s}} + s \varphi_{B(A_2 \ldots A_{2s}} \nabla_{A_1)B'} \xi^{BB'}
\]

\[
+ \frac{1}{2} \varphi_{A_1 \ldots A_{2s}} \nabla^{CC'} \xi_{CC'}.
\]

We shall now prove an auxiliary result on the derivatives of \( \eta_{AA'} \), which will allow us to prove our main result.

**Lemma 2.4.** Let \( \kappa_{AB} \in S_{2,0} \) satisfy the Killing spinor equation (2.2) and the aligned matter condition (2.4), and let \( \xi_{AA'} \) be given by (2.3). If \( \phi_{AB} \in S_{2,0} \) satisfies the Maxwell equation (2.4) and \( \eta_{AA'} \) is given by (2.8), then

\[
(\mathcal{P}_{1,1} \eta) = 0,
\]

\[
(\mathcal{G}_{1,1} \eta)_{AB} = \frac{2}{3} (\mathcal{L}_\xi \phi)_{AB},
\]

\[
(\mathcal{G}_{\dagger 1,1} \eta)_{A'B'} = 0,
\]

\[
\eta_{AA'} \xi^{AA'} = \kappa_{AB} (\mathcal{L}_\xi \phi)_{AB}.
\]
Proof of Theorem 1.1. From the Leibniz rule, we first find
\[ \nabla^{BB'} V_{AB'B'} = \frac{1}{2} \eta'_{ABB'} \nabla^{BB'} \eta'_{AB} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BA} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BB} + \frac{1}{2} \Theta_{AB} \nabla^{BB'} (\hat{\nabla} \phi)_{A'B'} + \frac{1}{2} \Theta_{A'B} \nabla^{BB'} (\hat{\nabla} \phi)_{AB}. \]
This can be simplified by first observing that \( \hat{\phi} \) is a symmetry operator taking solutions of the Maxwell equation to solutions of the Maxwell equation, so
\[ \nabla^{BB'} V_{AB'B'} = \frac{1}{2} \eta'_{ABB'} \nabla^{BB'} \eta'_{AB} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BA} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BB} + \frac{1}{2} \Theta_{AB} \nabla^{BB'} (\hat{\nabla} \phi)_{A'B'} + \frac{1}{2} \Theta_{A'B} \nabla^{BB'} (\hat{\nabla} \phi)_{AB}. \]
\[ \nabla^{BB'} V_{AB'B'} = \frac{1}{2} \eta'_{ABB'} \nabla^{BB'} \eta'_{AB} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BA} + \frac{1}{2} \eta_{A'B} \nabla^{BB'} \eta'_{BB} + \frac{1}{2} \Theta_{AB} \nabla^{BB'} (\hat{\nabla} \phi)_{A'B'} + \frac{1}{2} \Theta_{A'B} \nabla^{BB'} (\hat{\nabla} \phi)_{AB}. \]
\((\psi_2^\dagger \tilde{\zeta} \phi)_{AB} = 0\) and similarly for the complex conjugate. It can be further simplified by substituting the definition \(\nabla^{B}_{A} \Theta_{AB} = \eta_{AA'}\), cf. (2.14) to eliminate the derivative of \(\Theta_{AB}\) terms. This yields
\[
\nabla^{B'_{A'B'}} V_{ABA'B'} = - \frac{1}{2} \eta_{A'B'} (\psi_{1,1}\eta)_{AB} - \frac{1}{2} \eta^{B'}_{A'} (\bar{\psi}_{1,1}\bar{\eta})_{AB} - \frac{1}{2} \eta^{B'}_{A'} (\psi_{1,1}\eta)_{A'B'} - \frac{1}{2} \eta^{B'}_{A'} (\bar{\psi}_{1,1}\bar{\eta})_{A'B'} + \frac{1}{2} \eta_{A'B'} (\hat{\zeta} \otimes \bar{\phi})_{AB} + \frac{1}{2} \eta_{A'B'} (\hat{\bar{\zeta}} \otimes \phi)_{AB}.
\]
The terms involving \((\psi_{1,1}\eta)_{A'B'}\) and \((\bar{\psi}_{1,1}\bar{\eta})_{AB}\) are zero by equation (2.11). Those involving \((\hat{\bar{\eta}}_{1,1})\) and \((\psi_{1,1}\eta)_{AB}\) are zero by equation (2.11). Finally by equation (2.11), the terms involving \((\psi_{1,1}\eta)_{AB}\) and \((\bar{\psi}_{1,1}\bar{\eta})_{A'B'}\) cancel with those involving \((\hat{\zeta} \otimes \bar{\phi})_{AB}\) and \((\hat{\bar{\zeta}} \otimes \phi)_{A'B'}\) respectively. This completes the result. \(\square\)

3. Further remarks on the Kerr spacetime

The stationary, asymptotically flat, vacuum Kerr spacetimes, and more generally the electro-vacuum Kerr-Newman spacetimes, have algebraic type \(\{2, 2\}\), i.e. the Weyl spinor \(\Psi_{ABCD}\) has two distinct, repeated, principal spinors \(o_{A,A}\) which are unique up to a rescaling. The dyad \(o_{A,A}\) is normalized by \(o_{A,A}^2 = 1\). For the following discussion, recall that given a spin dyad \(o_{A,A}\), one defines for a symmetric spinor \(\tilde{\omega}_{A\cdots A}\), scalars \(\tilde{\omega}\) by contracting \(i\) times with \(\epsilon^A\) and \(k - i\) times with \(o^A\). This yields Weyl scalars \(\psi_i, i = 0, \ldots, 4\) and Maxwell scalars \(\phi_i, i = 0, 1, 2\).

In a spacetime of type \(\{2, 2\}\) with principal dyad \(o_{A,A}\), it holds that \(\Psi_{ABCD} = 6\Psi_{2,0\otimes 6,0\otimes 4,0}\), and in this case it follows from (2.11) that any valence \((2, 0)\) Killing spinor must be of the form
\[
\kappa_{AB} = \zeta^{(A'\Lambda)}_{\Lambda')(A'B),
\]
for some scalar \(\zeta\).

If \((t, r, \theta, \phi)\) are Boyer-Lindquist coordinates, then the Coulomb field, i.e. the unique static, regular Maxwell test field, on the Kerr-Newman spacetime takes the form
\[
\phi_{AB} = \frac{1}{(r - ia \cos \theta)^2} \theta_{(A'\Lambda)}^{(A'B)}
\]
up to a rescaling by a constant. In particular the extreme components \(\phi_0, \phi_2\) are zero. The background Maxwell field in the electro-vacuum Kerr-Newman spacetime is a constant multiple of this Coulomb field.

The Killing spinor \(\kappa_{AB}\) is
\[
\kappa_{AB} = \frac{2}{3} (r - ia \cos \theta) \theta_{(A'\Lambda)\Lambda(B)}
\]
which is therefore proportional to the background Maxwell field in the Kerr-Newman spacetime. Hence, by the Einstein equation, \(\Phi_{ABA'B'}\) is proportional to \(\kappa_{AB} \kappa_{A'B'}\). It follows that the aligned matter condition holds in the Kerr-Newman spacetime.

The normalisation in equation (3.2) is chosen so that \(\xi^a = (\partial_t)^a\), where \(\xi_a\) is given by (2.3). In particular \(\xi_a\) is real, which exhibits the fact that the Kerr-Newman family admits a Killing-Yano tensor, as remarked above. In particular, we see that the tensor \(V_{ab}\) given by (2.7) is conserved. More generally, any vacuum type \(\{2, 2\}\) spacetime admits a Killing spinor of valence \((2, 0)\), of the form (3.1) with \(\zeta\) proportional to \(\Psi_{2,1/3}\). This shows that Theorem 1.1 applies in the class of vacuum type \(\{2, 2\}\) metrics.
3.1. The Teukolsky equations and $V_{ab}$. The Maxwell equations on a Kerr black hole imply the $s = 1$ Teukolsky equations for the extreme scalars, $\phi_0$ and $\phi_2$. This system has many properties in common with the $s = 2$ Teukolsky equations which arise from linearising the Einstein equations. Despite the fact that the Teukolsky equations have been known for more than 40 years, and have been the subject of much study, no boundedness or decay estimates are known for the $s \neq 0$ Teukolsky equations, other than the mode stability result of Whiting [9].

In terms of the Maxwell scalars $\phi$, the Newman-Penrose scalars for $\Theta_{AB}$ satisfy
\begin{align*}
\Theta_0 &= -2\kappa_1 \phi_0, \\
\Theta_1 &= 0, \\
\Theta_2 &= 2\kappa_1 \phi_2.
\end{align*}
Thus, only the extreme components of $\phi_{AB}$ appear in $\Theta_{AB}$, and hence in $\eta_{AB}$. Equation (2.11d) can be used to express $(\hat{L}_\xi \phi)_{AB}$ in terms of $\eta_{AA'}$ and $(\hat{L}_\xi \Theta)_{AB}$, from which it follows that $V_{ab}$ can be written solely in terms of the extreme components of $\phi_{AB}$. This has two important consequences. Firstly, in the Kerr-Newman spacetime the extreme components of the Coulomb solutions vanish, and hence the conserved tensor $V_{ab}$ naturally excludes non-radiating solutions of the Maxwell equation. Secondly, since it is defined in terms of the extreme Maxwell scalars alone, $V_{ab}$ can be thought of as an “energy-momentum tensor” for the $s = 1$ combined Teukolsky/Teukolsky-Starobinsky system, which corresponds to equations (2.11b)-(2.11c), see [3] for more details.

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