On the spherically symmetric Einstein–Yang–Mills–Higgs equations in Bondi coordinates

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We revisit and generalize, to the Einstein–Yang–Mills–Higgs (EYMH) system, previous results of Christodoulou and Chae concerning global solutions for the Einstein-scalar field and the Einstein–Maxwell–Higgs (EMH) equations. The novelty of the present work is twofold. For one thing, the assumption on the self-interaction potential is improved. For another thing, explanation is furnished why the solutions obtained here and those proved by Chae for the EMH system decay more slowly than those established by Christodoulou in the case of self-gravitating scalar fields. Actually, this latter phenomenon stems from the non-vanishing local charge in EMH and EYMH models.

Keywords: global solution; Einstein–Yang–Mills–Higgs equations; spherical symmetry; Bondi coordinates

1. Introduction

The Yang–Mills–Higgs (YMH) field equations arise in elementary particle physics. The Yang–Mills equations appear as the generalization of the classical Maxwell equations where ordinary derivatives are replaced by covariant derivatives. The YMH equations are nonlinear partial differential equations (PDEs) that are conformally invariant and gauge-invariant. This latter property has been exploited by Eardley & Moncrief (1982a,b) to make a remarkable contribution to the area of PDE in proving global existence for these equations in four-dimensional Minkowski space.

When the YMH equations are coupled to gravity, such a global existence result is yet to be proved without any symmetry assumption. Since this problem is a very difficult one, it makes sense to start by investigating it with some symmetry assumption. In this paper, we prove a global existence result for the Einstein–Yang–Mills–Higgs (EYMH) equations under the assumption of spherical symmetry.

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The issue of proving global existence results for gravitational and matter field equations is of interest in mathematical relativity and in the area of PDE. In general relativity, the global existence problem is important since it is a reformulation of the cosmic censorship conjecture rendering this more amenable to direct analytic attack (Eardley & Moncrief 1981).

Much work has been done in the past concerning the EYMH equations. Let us mention for instance that the spontaneous compactification of space in EYMH model has been studied in Cremer & Schrek (1976). A canonical formulation of the spherically symmetric EYMH system for a general gauge group has been found in Cordero (1977). Regular localized solutions of the electric or magnetic type were found in Clément (1981). Stimulated by the discovery of globally regular solutions of the EYM equations, numerically by Bartnik & McKinnon (1988), many authors published a considerable number of papers dealing with the static spherically or axially symmetric EYMH equations. Most of these papers concern asymptotically flat, smooth particle-like and black hole solutions (Hajicek 1983a, b; Winstanley & Mavromaos 1995; Gal’tsov & Donets 1997; Brihaye & Desoil 2000; Kleihaus & Kunz 2000; Hartmann et al. 2001; Paturyan et al. 2005) although some cosmological solutions have been found for instance in Breitenlohner et al. (2000) and Forgács & Reuillon (2005).

As far as rigorous existence proofs are concerned, Künzle & Oliynyk (2006) used differential geometry techniques to investigate the static spherically symmetric EYMH system. Local existence of the characteristic Cauchy problem for EYMH without any symmetry assumption has been proved recently by Dossa & Tadmon (2010a, b). More recently, Vacaru (2011) has written the Einstein field equations in variables adapted to non-holonomic $2 + 2$ splitting and has applied this geometric technique for constructing off-diagonal exact solutions of EYMH equations.

In the present paper, as mentioned already, we investigate the initial value problem for the spherically symmetric EYMH equations. We write the system of equations in the so-called Bondi coordinates so as to obtain a characteristic Cauchy problem, then we prove global existence and some decay property of solutions. This setting has been used before by Christodoulou (1986) for the Einstein-scalar field system and by Chae (2001) and Chae (2003), respectively, for Einstein–Klein–Gordon (EKG) and Einstein–Maxwell–Higgs (EMH) equations. To reach our goal, we take advantage of the tools set up in Christodoulou (1986) to reduce the problem to that of the resolution of a nonlinear evolution system of two PDE. By comparison with the works in Chae (2003) and Christodoulou (1986), where a single evolution equation was dealt with, supplementary difficulties arise: here the number of terms to estimate is much more higher and some additional terms have to be handled meticulously. By dint of arduous calculations combined with several handy mathematical tools, we implement a fixed point method to arrive, under a more general assumption on the self-interaction potential, at a global existence and uniqueness result for the spherically symmetric EYMH system with appropriate initial data. Furthermore, we show that this solution decays more slowly than the one obtained in Christodoulou (1986). A thorough examination reveals that this latter phenomenon stems from the presence of the local charge $Q$. Some questions raised by Chae (2003) are thus answered. It could be interesting to see whether the decay properties of the solution mentioned earlier can be explained by using compactification techniques due to Penrose (1964). In final, we obtained a generalization and improvement of the results in Chae (2003).
The spherically symmetric EYMH equations

and Christodoulou (1986). The main result of the present paper encompasses the EYMH system with zero self-interaction potential and the EKG system as well. It is worth noting that this paper is the corrected and detailed version of the Note Tadmon (2011).

The present work is organized as follows. Section 2 is devoted to the derivation of EYMH equations under the spherical symmetry assumption. In §3, we show that the equations reduce to a nonlinear evolution system. In §4, we state and prove the main result of the present investigation.

2. The Einstein–Yang–Mills–Higgs system

(a) Equations of motion

Throughout the paper, unless stated otherwise, Einstein convention is used, e.g. \( w^d_c h^c = \sum_e w^d_e h^e \). We concentrate on the \( \mathfrak{su}(2) \)-EYMH and assume that the Yang–Mills field is in the adjoint representation of \( \mathfrak{su}(2) \) while the Higgs fields is in the fundamental representation of \( \mathfrak{su}(2) \). The basic elements of the \( \mathfrak{su}(2) \)-EYMH system consist of a quadruplet \((\mathcal{M}, g, A, \Phi)\), where \( \mathcal{M} \) is a 4D space–time manifold equipped with a Lorentzian metric \( g \); \( A \) is a 1–form, called the Yang–Mills potential, defined on \( \mathcal{M} \) with values in the Lie algebra \( \mathfrak{su}(2) \) of the Lie group \( SU(2) \); \( \Phi \) is a scalar-multiplet field, called the Higgs field, defined on \( \mathcal{M} \). The fields equations for the \( \mathfrak{su}(2) \)-EYMH model read (Choquet-Bruhat 1991; Dossa & Tadmon 2010a, b; Tadmon 2011)

\[
\begin{align*}
R_{\alpha\beta} - \frac{1}{2} g_{\alpha\beta} R &= T_{\alpha\beta}, \\
g^{\lambda\mu}(\nabla_\lambda F_{\mu\alpha} + [A_\lambda, F_{\mu\alpha}]) &= J_\alpha,
\end{align*}
\]

and

\[
\begin{align*}
g^{\lambda\mu} \hat{\nabla}_\lambda \hat{\nabla}_\mu \Phi &= V'(\Phi^\dagger \Phi) \Phi.
\end{align*}
\]

Here, \( g_{\alpha\beta} \) are the components of the space–time metric. \( R_{\alpha\beta} \) and \( R \) are, respectively, the Ricci tensor and the scalar curvature relative to the space–time metric. \( (F_{\alpha\beta}) \) represents the Yang–Mills strength field \( F \), which is a \( \mathfrak{su}(2) \)-valued antisymmetric 2-form of type \( Ad \), defined on \( \mathcal{M} \). \( F \) is related to the unknown Yang–Mills potential \( A \) as follows:

\[
F_{\alpha\beta} = \nabla_\alpha A_\beta - \nabla_\beta A_\alpha + [A_\alpha, A_\beta].
\]

\( \nabla \) denotes the covariant derivative w.r.t. the space–time metric \( g \); \([,] \) denote the Lie brackets of the Lie algebra \( \mathfrak{su}(2) \). \( \hat{\nabla}_\alpha \Phi \) is the gauge covariant derivative of the complex doublet Higgs field, defined by

\[
\hat{\nabla}_\alpha \Phi = \frac{\partial \Phi}{\partial x^\alpha} - iA_\alpha \frac{\sigma_I}{2} \Phi,
\]

where \( (\sigma_I)_{I=1,2,3} \) are the conventional Pauli spin matrices. \( \Phi^\dagger \) denotes the hermitian conjugate of \( \Phi \), i.e. \( \Phi^\dagger \) is the transpose of \( \Phi^* \), where \( \Phi^* \) is the complex conjugate of \( \Phi \). \( V \) is a real function defined on \([0, \infty)\), often called the self-interaction potential, with derivative \( V' \). \( T_{\alpha\beta} \) are the components of the
energy–momentum–stress tensor, given by
\[ T_{\alpha\beta} = g^{\rho\sigma} F_{\alpha\rho} F_{\beta\sigma} - \frac{1}{4} g_{\alpha\beta} g^{\sigma\lambda} g^{\rho\mu} F_{\sigma\rho} F_{\lambda\mu} \]
\[ + (\hat{\nabla}_\alpha \Phi)^\dagger \hat{\nabla}_\beta \Phi + (\hat{\nabla}_\beta \Phi)^\dagger \hat{\nabla}_\alpha \Phi - g_{\alpha\beta} (g^{\sigma\rho} (\hat{\nabla}_\sigma \Phi)^\dagger \hat{\nabla}_\rho \Phi + V(\Phi^\dagger \Phi)), \]
where the dot ‘.’ denotes the \( Ad \)-invariant non-degenerate scalar product of \( su(2) \) defined by
\[ f, k = \sum_{I=1}^{3} f^I k^I, \quad f \in su(2), \quad k \in su(2). \] (2.5)
It is worth noting that the \( Ad \)-invariant non-degenerate scalar product ‘.’ defined in (2.5) enjoys the following property:
\[ f, [k, l] = [f, k], l, \quad \forall f, k, l \in su(2). \] (2.6)
\( J_\alpha \) are the components of the Yang–Mills current, given by
\[ J_\alpha^I = \Phi^I S^I \hat{\nabla}_\alpha \Phi - (\hat{\nabla}_\alpha \Phi)^\dagger S^I \Phi, \quad I = 1, 2, 3, \quad \alpha = 0, 1, 2, 3; \] (2.7)
where
\[ S^I = i \frac{\sigma^I}{2}. \]
We work within the basis \((T_I)_{I=1,...,3}\) of \( su(2) \), with
\[ T_1 = i \frac{\sigma_1}{2}, \quad T_2 = -i \frac{\sigma_2}{2} \quad \text{and} \quad T_3 = i \frac{\sigma_3}{2}. \]
(b) The spherically symmetric ansätze and fields equations

We will work in a Bondi coordinates system \((x^\alpha) = (u, r, \theta, \phi)\) on \( \mathbb{R}^4 \), used in series of works by Christodoulou (1986) and Chae (2001, 2003), where \( u \) is a retarded time coordinate and \( r \) is a radial coordinate. In this coordinates system, the most general form for the spherically symmetric metric could be written as follows:
\[ ds^2 = -e^{2\nu} du^2 - 2e^{\nu+\lambda} du dr + r^2 (d\theta^2 + \sin^2 \theta d\phi^2), \] (2.8)
where \( \nu \) and \( \lambda \) are real functions of \( u \) and \( r \) only.

The general form for the spherically symmetric Yang–Mills potential in the adjoint representation of \( su(2) \) could be written as follows (Choquet-Bruhat 1991; Dossa & Tadmon 2010a,b):
\[ A = a T_3 du, \] (2.9)
i.e. \( A_0 = a T_3, \quad A_1 = A_2 = A_3 = 0 \), where \( a \) is a function of \( u \) and \( r \) only.

We will use the following ansatz for the spherically symmetric Higgs field in the fundamental representation of \( su(2) \)
\[ \Phi = \begin{pmatrix} 0 \\ \psi + i\xi \end{pmatrix}, \] (2.10)
where \( \psi \) and \( \xi \) are real functions of \( u \) and \( r \) only.

As noted in Tadmon & Tchapnda (2012), writing roughly the EYMH fields equations in the coordinates system \((u, r, \theta, \phi)\) yields a system of PDE that is very difficult to handle. Ahead of overcoming this toughness, we introduce,
as in Christodoulou (1986), a null tetrad \( (e_\alpha)_{\alpha = 0, \ldots, 3} \) defined by \( e_0 = e^{-\nu}(\partial/\partial u) - (1/2)e^{-\lambda}(\partial/\partial r) \), \( e_1 = e^{-\lambda}(\partial/\partial u) \) and \( (e_2, e_3) \) is a locally defined orthonormal frame on the unit 2-sphere. Denote by over dot (\( . \)) differentiation with respect to \( u \), and by a prime (\( ' \)) differentiation with respect to \( r \), e.g. \( \dot{\lambda} = \partial\lambda/\partial u \), \( \nu' = \partial\nu/\partial r \). After many lengthy calculations (Tadmon & Tchapnda 2012), the relevant Einstein fields equations \( R(e_\alpha, e_\beta) = T(e_\alpha, e_\beta) \) are the following:

\[
\frac{1}{2}(\lambda' + \nu') - 2e^{\lambda-\nu} \dot{\lambda} = 2re^{2(\lambda-\nu)}[(\dot{\psi} - \frac{1}{2}a\xi)^2 + (\dot{\xi} + \frac{1}{2}a\psi)^2] \\
- 2re^{\lambda-\nu}[(\dot{\psi} - \frac{1}{2}a\xi)\psi' + (\dot{\xi} + \frac{1}{2}a\psi)\xi'] \\
+ \frac{1}{2}r[\psi'^2 + (\xi')^2],
\]

(2.11)

\[
\nu' - \lambda' + r^{-1}(1 - e^{2\lambda}) + r\left(\frac{1}{2}e^{-2\nu}(a')^2 + e^{2\lambda} V\right) = 0,
\]

(2.12)

\[
\lambda' + \nu' = r[(\psi')^2 + (\xi')^2]
\]

(2.13)

and

\[
\nu'' + (\nu' - \lambda')(\nu' + r^{-1}) - e^{\lambda-\nu}[(\dot{\lambda})' + (\dot{\nu})' + 2(\dot{\psi} - \frac{1}{2}a\xi)\psi' \\
+ 2(\dot{\xi} + \frac{1}{2}a\psi)\xi'] - \frac{1}{2}e^{-2\nu}(a')^2 + [(\psi')^2 + (\xi')^2] + e^{2\lambda} V = 0.
\]

(2.14)

The relevant Yang–Mills equations are the following:

\[
a'' + (2r^{-1} - \lambda' - \nu')a' - e^{\lambda-\nu}((\dot{a})' - \dot{\lambda} - \dot{\nu}) = e^{2\lambda}[(\dot{\psi} - \dot{\xi} + \frac{1}{2}a(\psi^2 + \xi^2) \]

(2.15)

and

\[
a'' + (2r^{-1} - \lambda' - \nu')a' = e^{\lambda+\nu}(\psi\xi' - \xi\psi').
\]

(2.16)

The Higgs equations are found to be equivalent to the following system:

\[
\begin{aligned}
2(\dot{\psi}') - a\xi' - \frac{1}{2}a'\xi + 2r^{-1}(\dot{\psi} - \frac{1}{2}a\xi) \\
- e^{\nu-\lambda}[\psi'' + \psi'(2r^{-1} + \nu' - \lambda')] &= -\psi V'e^{\lambda+\nu}
\end{aligned}
\]

(2.17)

and

\[
\begin{aligned}
2(\dot{\xi}') + a\psi' + \frac{1}{2}a'\psi + 2r^{-1}(\dot{\xi} + \frac{1}{2}a\psi) \\
- e^{\nu-\lambda}[\xi'' + \xi'(2r^{-1} + \nu' - \lambda')] &= -\xi V'e^{\lambda+\nu}.
\end{aligned}
\]

Remark 2.1.

(i) It is shown in Tadmon (2012) that if the YMH equations are satisfied, then the energy–momentum–stress tensor given by (2.4) is divergence-free, i.e. \( g^{\alpha\sigma}\nabla_\sigma T_{\alpha\beta} = 0 \). Setting, as in Christodoulou (1986), \( E_{\alpha\beta} = R_{\alpha\beta} - (1/2)Rg_{\alpha\beta} - T_{\alpha\beta,} \) and using the contracted Bianchi identity \( g^{\alpha\sigma}\nabla_\sigma(R_{\alpha\beta} - (1/2)Rg_{\alpha\beta}) = 0 \), it easily follows that:

\[
g^{\alpha\sigma}\nabla_\sigma E_{\alpha\beta} = 0,
\]

(2.18)

provided that the YMH equations are satisfied.

(ii) It is straightforward to check by using (2.18) that, if the Einstein equations (2.12) and (2.13), i.e. \( E(e_1, e_1) = 0 \) and \( E(e_0, e_1) = 0 \), are satisfied as well as the YMH equations (2.15), (2.16) and (2.17), then the Einstein equation (2.14) is satisfied, i.e. \( E(e_2, e_2) = E(e_3, e_3) = 0 \). Furthermore,
\[ E(e_0, e_0) \] solves the following PDE:

\[ [E(e_0, e_0)]' + 2(v' + r^{-1})E(e_0, e_0) = 0, \]

which yields

\[ E(e_0, e_0)(u, r) = \left( \frac{r_0}{r} \right)^2 E(e_0, e_0)(u, r_0) \exp[2(v(u, r_0) - v(u, r))]. \]

This implies \( E(e_0, e_0) = 0 \) if regularity is assumed at the centre.

(iii) From the earlier-mentioned observations, we conclude that the Einstein equations \( (2.12) \) and \( (2.13) \), together with the YMH equations \( (2.15)-(2.17) \), are equivalent to the full set of EYMH equations \( (2.11)-(2.17) \).

3. Reduction of the Einstein–Yang–Mills–Higgs equations to a nonlinear evolution system

We adapt the tools set up and implemented in Christodoulou (1986). Two new functions are introduced as follows:

\[ h = (r \psi)' \quad \text{and} \quad k = (r \xi)'. \]

Then

\[ \psi = h = \frac{1}{r} \int_0^r h(s) \, ds, \quad \xi = k = \frac{1}{r} \int_0^r k(s) \, ds, \quad \psi' = \frac{h - \bar{h}}{r} \quad \text{and} \quad \xi' = \frac{k - \bar{k}}{r}. \]

The Einstein equation \( (2.13) \) then reads

\[ \lambda' + \nu' = \frac{1}{r}[(h - \bar{h})^2 + (k - \bar{k})^2], \]

and the solution, which satisfies the asymptotic condition \( \lambda + \nu \to 0 \) as \( r \to \infty \), is

\[ \lambda + \nu = -\int_r^{+\infty} \frac{1}{s}[(h - \bar{h})^2 + (k - \bar{k})^2] \, ds. \] (3.3)

Integrating the Yang–Mills equation \( (2.16) \) yields

\[ a(u, r) = \int_0^r e^{\lambda + \nu} \frac{Q(u, s)}{s^2} \, ds, \] (3.4)

where \( Q \) is the local charge function defined by

\[ Q(u, r) = \int_0^r s(hk - \bar{k}h)(u, s) \, ds. \] (3.5)

Now we put the Einstein equation \( (2.12) \) under the form

\[ (v' - \lambda')e^{v-\lambda} + r^{-1}e^{v-\lambda} = r^{-1}e^{\lambda + \nu} - r\left(\frac{1}{2}e^{-\lambda - v}(a')^2 + e^{\lambda + \nu} V\right), \]

which is integrated to give

\[ e^{v-\lambda} = \frac{1}{r} \int_0^r e^{\lambda + \nu} \left[ 1 - \frac{1}{2} \frac{Q^2}{s^2} - s^2 V \right] \, ds. \] (3.6)
From (2.17) and (3.1), we recast the Higgs equations into the following nonlinear evolution system:

\[
\begin{aligned}
\dot{h} - \frac{1}{2} e^{\nu - \lambda} h' &= \frac{1}{2r} (e^{\lambda + \nu} - e^{\nu - \lambda}) (h - \bar{h}) - \frac{Q^2}{4r^3} e^{\lambda + \nu} (h - \bar{h}) \\
&\quad + \frac{Q}{4r} \tilde{h} e^{\lambda + \nu} - \frac{1}{2} [r (h - \bar{h}) V e^{\lambda + \nu} + r \bar{h} V' e^{\lambda + \nu} - ak], \\
\end{aligned}
\]

and

\[
\begin{aligned}
\dot{k} - \frac{1}{2} e^{\nu - \lambda} k' &= \frac{1}{2r} (e^{\lambda + \nu} - e^{\nu - \lambda}) (k - \bar{k}) - \frac{Q^2}{4r^3} e^{\lambda + \nu} (k - \bar{k}) \\
&\quad + \frac{Q}{4r} \tilde{k} e^{\lambda + \nu} - \frac{1}{2} [r (k - \bar{k}) V e^{\lambda + \nu} + r \bar{k} V' e^{\lambda + \nu} + ah]. \\
\end{aligned}
\]

We write system (3.7) in matrix form with unknown vector function \( W = (h, k) \) as follows:

\[
DW = \frac{1}{2r} (g - \tilde{g})(W - \tilde{W}) - \frac{Q^2 g}{4r^3} (W - \tilde{W}) + \frac{Qg}{4r} \sigma_1 \tilde{W} + \frac{a}{2} i\sigma_2 W \\
- \frac{r}{2} [Vg(W - \tilde{W}) + V'g\tilde{W}],
\]

where

\[
D = \frac{\partial}{\partial u} - \frac{\tilde{g}}{2} \frac{\partial}{\partial r}, \quad g = e^{\lambda + \nu}, \quad \tilde{g} = e^{\nu - \lambda}.
\]

**Remark 3.1.**

(i) It is worth emphasizing that the integration of the Einstein equations is achieved under the asymptotic condition \( \lambda + \nu \to 0, \ r \to \infty \). The Yang–Mills equation (2.16) is integrated easily by classical tools (change of unknown function and variation of the constant) to yield the solution \( a(u, r) \) that exists for all \( r \in [0, +\infty) \).

(ii) For the Yang–Mills field, it is easy to see that it vanishes at spatial infinity i.e. \( F(u, r) \to 0, \ r \to +\infty \). Actually, by simple calculation, all the components of the Yang–Mills field \( F \) vanish except \( F_{01} \), which is given by \( F_{01} = -a'T_3 \). Using the expression of \( a(u, r) \) and estimating the local charge \( Q(u, r) \), one easily gains \( F_{01}(u, r) \to 0, \ r \to +\infty \).

(iii) It could be of interest to investigate the case of an asymptotic (anti) de Sitter metric.

4. Existence and uniqueness of global classical solutions

This section aims at stating and proving a global existence and uniqueness result for the initial value problem associated with the nonlinear evolution system (3.8). In order to do this, we begin by introducing the spaces of functions used and some preliminary notations.
(a) Functional framework, notations and statement of the main result

For a vector function, \( W = (\frac{h}{k}) \) defined on \([0, \infty) \times [0, \infty)\) with the real functions \( h \) and \( k \) belonging both to \( C^1([0, \infty) \times [0, \infty)) \), we will just write \( W \in C^1([0, \infty) \times [0, \infty)) \) instead of \( W \in [C^1([0, \infty) \times [0, \infty))]^2 \). We will also use the notation \( |W| := |h| + |k| \). Consider the initial value problem for system (3.8) with initial datum \( W_0(r) = W(0, r) \). Following the works in Chae (2003), we define the Banach function spaces \((\mathcal{X}, \|\cdot\|_\mathcal{X}), (\mathcal{X}_0, \|\cdot\|_{\mathcal{X}_0})\) and \((\mathcal{Y}, \|\cdot\|_\mathcal{Y})\) by

\[
\mathcal{X} = \{ W = W(u, r) \in C^1([0, \infty) \times [0, \infty)) : \|W\|_\mathcal{X} < \infty \},
\]

\[
\mathcal{X}_0 = \{ v = v(r) \in C^1([0, \infty)) : \|v\|_{\mathcal{X}_0} < \infty \}
\]

and

\[
\mathcal{Y} = \{ W = W(u, r) \in C^1([0, \infty) \times [0, \infty)) : W(0, r) = W_0(r), \quad \|W\|_\mathcal{Y} < \infty \},
\]

where

\[
\|W\|_\mathcal{X} := \sup_{u, r \geq 0} \{(1 + r + u)^2|W(u, r)| + (1 + r + u)^3|W'(u, r)|\},
\]

\[
\|v\|_{\mathcal{X}_0} := \sup_{r \geq 0} \{(1 + r)^2|v(r)| + (1 + r)^3|v'(r)|\}
\]

and

\[
\|W\|_\mathcal{Y} := \sup_{u, r \geq 0} \{(1 + r + u)^2|W(u, r)|\}.
\]

We are now in the position to state the main result of this paper, which is the generalization, as we have mentioned from the outset, of the results obtained in Christodoulou (1986) and Chae (2003).

**Theorem 4.1.** Assume for the self-interaction potential \( V \) that \( V \in C^2([0, \infty)) \), and there exists a constant \( K_0 \geq 0 \) such that

\[
|V(t)| + t|V'(t)| + t^2|V''(t)| \leq K_0 t^{p+1} \quad \forall t \geq 0, \text{ for some } p \geq \frac{3}{2}.
\]

Suppose for the initial datum \( W_0 \) that

\[
W_0 \in C^1([0, \infty)), \quad W_0(r) = O(r^{-2}) \quad \text{and} \quad W_0'(r) = O(r^{-3}).
\]

Then there exists \( \epsilon > 0 \) such that if \( \|W_0\|_{\mathcal{X}_0} < \epsilon \), then there exists a unique global classical solution \( W \in C^1([0, \infty) \times [0, \infty)) \) of (3.8) such that \( W(0, r) = W_0(r) \). In addition, this solution fulfills the decay property

\[
|W(u, r)| \leq C(1 + u + r)^{-2} \quad \text{and} \quad |W'(u, r)| \leq C(1 + u + r)^{-3},
\]

where \( C > 0 \) is an increasing continuous function of \( K_0 \). Moreover, the corresponding space–time is time-like and null geodesically complete towards the future.

**Remark 4.2.**

(i) Theorem 4.1 was stated and proved, under more restrictive assumptions, by Chae (2003) for the EMH model. Notably, in Chae (2003), the assumption on \( p \) was \( p \geq 3 \). It is our aim, in this paper, to provide a generalization and improvement of the results obtained in Christodoulou (1986) and Chae (2003) by extending them, under assumption (4.2),

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to the EYMH model. Note also that, unlike the solution obtained by Christodoulou (1986) for the spherically symmetric Einstein-scalar field system, the solution obtained in the present work decays more slowly than that of Christodoulou (1986). This latter fact is attributed, as it will be shown thereafter, to the non-vanishing of the local charge $Q$.

(ii) Note that assumption (4.2) is satisfied for $V \equiv 0$. Theorem 4.1 then applies to provide global existence and uniqueness of classical solutions of the spherically symmetric EYMH system with vanishing self-interaction potential.

(iii) If $\Phi$ is a real doublet of the form $\Phi = (0 \iota)$, i.e. $\xi \equiv 0$, then the local charge $Q$ as well as the Yang–Mills potential $A$ and the Yang–Mills strength field $F$ vanish (see (3.1), (3.5), (3.4) and (2.2)). In this case, we can choose $V(t) = t^{p+1}/(p+1)$, $t \geq 0$, so that the Higgs equations reads

\begin{equation}
\psi^{2p+1}.
\end{equation}

(4.5) is the nonlinear Klein–Gordon equation that can be reduced to the form

\begin{equation}
Dh = \frac{1}{2r}(g - \tilde{g})(h - \tilde{h}) - \frac{rg}{2} \left[ (h - \tilde{h}) (\tilde{h})^{2p+2} + (\tilde{h})^{2p+1} \right],
\end{equation}

where

\begin{equation}
\tilde{g} = \frac{1}{r} \int_{0}^{r} g \left[ 1 - s^2 \frac{(\tilde{h})^{2p+2}}{p+1} \right] ds.
\end{equation}

Theorem 4.1 then applies in this situation to encompass global existence and uniqueness of classical spherically symmetric solutions of the nonlinear EKG system. Moreover, in this case, it turns out that the solutions obtained possess the same order of decay estimates as those obtained by Christodoulou (1986). The same statement holds true if $\Phi$ is of the form $\Phi = (0 \iota)$, i.e. $\psi \equiv 0$.

(b) Proof of theorem 4.1

It is worth noting that some computational details can be found in Tadmon & Tchapnda (2012). Throughout this paragraph, given that \{I_2, \sigma_1, i\sigma_2, \sigma_3\} is a basis for the real vector space of $2 \times 2$ real matrices, $I_2$ being the $2 \times 2$ identity matrix, we will use the notation

\begin{equation}
|N| = |n_0| + |n_1| + |n_2| + |n_3|,
\end{equation}

if $N$ is a $2 \times 2$ real matrix with

\begin{equation}
N = n_0 I_2 + n_1 \sigma_1 + n_2 (i\sigma_2) + n_3 \sigma_3.
\end{equation}

We prove theorem 4.1 by a contraction mapping argument as in Chae (2003), Christodoulou (1986), Tadmon (2011) while improving and correcting at the same time some key estimates therein. We define the mapping $\mathcal{K} : W \rightarrow w = \mathcal{K}(W)$,
where \( w \) is the solution of the first-order linear initial value problem

\[
Dw = \frac{1}{2r}(g - \tilde{g})(w - \bar{W}) - \frac{Q^2 g}{4r^3} (w - \bar{W}) + \frac{Qg}{4r} \sigma_1 \bar{W} \\
+ \frac{a}{2} i \sigma_2 w - \frac{r}{2} [Vg(w - \bar{W}) + V'g \bar{W}]
\]

(4.6)

and

\[
w(0, r) = W_0(r).
\]

Our purpose is to show that the mapping \( K \) defined earlier is a contraction from a non-empty closed ball of \( X \) into itself. If this is done, then the standard fixed point theorem applies to yield the unique fixed point \( W \in X \) such that \( K(W) = W \), which is the solution of the nonlinear system (3.8) satisfying \( W(0, r) = W_0(r) \).

(i) \( K \) is a mapping from a ball of \( X \) into itself

Let \( \rho > 0 \) be a real number. \( B_\rho \) denotes the closed ball, in \( X \), of radius \( \rho \) centred at 0, i.e.

\[
B_\rho = \{ W \in X : \| W \|_X \leq \rho \}.
\]

We will prove that \( \rho \) can be chosen small enough such that \( K : B_\rho \to B_\rho \). Let \( W \in B_\rho \). We have to estimate \( \| K(W) \|_X \) in terms of \( \| W \|_X \). The characteristic system of ODE associated with the initial value problem (4.6) is

\[
\begin{align*}
\frac{dr}{du} &= -\frac{1}{2} \tilde{g} \\
\frac{dw}{du} &= \frac{1}{2r}(g - \tilde{g})(w - \bar{W}) - \frac{Q^2 g}{4r^3} (w - \bar{W}) + \frac{Qg}{4r} \sigma_1 \bar{W} \\
&\quad + \frac{a}{2} i \sigma_2 w - \frac{r}{2} [Vg(w - \bar{W}) + V'g \bar{W}],
\end{align*}
\]

(4.7)

with initial data

\[
r(0) = r_0 \quad \text{and} \quad w(0) = W_0.
\]

Let \( r(u) = \gamma(u, r_0) \) be the solution of the initial value problem

\[
\frac{dr}{du} = -\frac{1}{2} \tilde{g}(u, r), \quad r(0) = r_0.
\]

(4.8)

Then

\[
r_1 = r_0 - \frac{1}{2} \int_0^{u_1} \tilde{g}(u, \gamma(u, r_0)) du,
\]

(4.9)

where \( r_1 = \gamma(u_1, r_0) \). Integrating the second ODE of (4.7) along \( \gamma \), we obtain

\[
w(u_1, r_1) = \exp \left( \int_0^{u_1} [N(u, r)]_{\gamma} dv \right) W(0, r_0) \\
+ \int_0^{u_1} \left\{ \exp \left( \int_u^{u_1} [N(u, r)]_{\gamma} dv \right) \right\} [f]_{\gamma} du,
\]

(4.10)
where the matrix function $N$ and the vector function $f$ are given by

$$N(u, r) = \left( \frac{1}{2r} (g - \tilde{g}) - \frac{Q^2 g}{4r^3} - \frac{r}{2} Vg \right) I_2 + \frac{a}{2} i \sigma_2$$

and

$$f(u, r) = \left( \frac{Q^2 g}{4r^3} - \frac{1}{2r} (g - \tilde{g}) + \frac{rg}{2} [V - V'] \right) \tilde{W} + \frac{Qg}{4r} \sigma_1 \dot{W}.$$  \hfill (4.11)

We have to estimate $N$ and $f$. Setting $\|W\|_X = x$, we derive the following estimates as in Chae (2003):

$$|\sigma_1 \tilde{W}(u, r)| = |\tilde{W}(u, r)| \leq \frac{x}{(1 + u)(1 + u + r)}$$  \hfill (4.12)

and

$$|W(u, r) - \tilde{W}(u, r)| \leq \frac{xr}{2(1 + u)(1 + u + r)^2}.$$  \hfill (4.13)

(4.13) implies

$$\int_0^\infty \frac{|W(u, r) - \tilde{W}(u, r)|^2}{r} \, dr \leq \frac{x^2}{24(1 + u)^4}.$$  \hfill (4.14)

In view of (3.3) and the definition of $g$ in (3.9), (4.14) yields

$$g(u, 0) \geq \exp \left( -\frac{x^2}{24(1 + u)^4} \right).$$  \hfill (4.15)

From (4.13), we gain

$$|g(u, r) - \tilde{g}(u, r)| \leq \frac{x^2 r^2}{12(1 + u)^3(1 + u + r)^3}.$$  \hfill (4.16)

In view of (3.5) and (4.12), we estimate the local charge to get

$$|Q(u, r)| \leq \frac{x^2 r^2}{2(1 + u)^2(1 + u + r)^2}.$$  \hfill (4.17)

Hence

$$\frac{1}{r} \int_0^r \frac{|Q(u, s)|^2}{s^2} \, ds \leq \frac{x^4 r^2}{12(1 + u)^5(1 + u + r)^3}.$$  \hfill (4.18)

The term containing the self-interaction potential is estimated using $\Phi^i \Phi = (\tilde{h})^2 + (\tilde{k})^2$ and assumption (4.2) as

$$\left| \frac{1}{r} \int_0^r s^2 V(\Phi^i \Phi) \, ds \right| \leq \frac{K_0 x^{2p+2} r^2}{3(1 + u)^6(1 + u + r)^4}.$$  \hfill (4.19)

From the definition of $\tilde{g}$ in (3.9) and the estimates (4.16), (4.18) and (4.19), we get

$$|g(u, r) - \tilde{g}(u, r)| \leq \frac{(3 + 8K_0)(x^2 + x^4 + x^{2p+2}) r^2}{24(1 + u)^5(1 + u + r)^3}.$$  \hfill (4.20)
It is easy to see that the function
\[ f \geq \exp \left( -\frac{x^2}{24} \right) - \frac{x^4}{24} - \frac{K_0 x^{2p+2}}{3}. \]  
(4.21)

It is easy to see that the function \( l = l(x) = \exp(-x^2/24) - \frac{x^4}{24} - \frac{K_0 x^{2p+2}}{3} \) has

a unique positive root \( x_0 \) and \( l(x) \in (0, 1) \) for all \( x \in [0, x_0) \). We now estimate

the function \( f \) defined in (4.11) as follows:

\[ |f(u, r)| \leq \left| \frac{Q^2 g \bar{W}}{4r^3} \right| + \left| \frac{1}{2r} (g - \bar{g}) \tilde{W} \right| + \left| \frac{rg}{2} (V - V') \tilde{W} \right| + \left| \frac{Qg}{4r} \sigma_1 \tilde{W} \right|. \]  
(4.22)

Since \( 0 < g \leq 1 \), (4.12) and (4.17) give

\[ \left| \frac{Q^2 g \bar{W}}{4r^3} \right| \leq \frac{x^5 r}{16(1 + u)^5(1 + u + r)^5} \]  
(4.23)

and

\[ \left| \frac{Qg}{4r^3} \sigma_1 \tilde{W} \right| \leq \frac{x^3 r}{8(1 + u)^3(1 + u + r)^3}. \]  
(4.24)

Using (4.12) and (4.20), we gain

\[ \left| \frac{1}{2r} (g - \bar{g}) \tilde{W} \right| \leq \frac{(3 + 8K_0)(x^3 + x^5 + x^{2p+3}) r}{48(1 + u)^6(1 + u + r)^4}. \]  
(4.25)

Assumption (4.2) and the estimate (4.12) imply

\[ \left| \frac{rg}{2} (V(\Phi \Phi') - V'(\Phi \Phi)) \tilde{W} \right| \leq \frac{K_0}{2} \left[ \frac{x^{2p+3}}{(1 + u)^6(1 + u + r)^5} + \frac{x^{2p+1}}{(1 + u)^4(1 + u + r)^3} \right] \]  
(4.26)

Inserting (4.23)–(4.26) into (4.22) yields

\[ |f(u, r)| \leq \frac{(12 + 56K_0)(x^3 + x^5 + x^{2p+1} + x^{2p+3})}{48(1 + u)^3(1 + u + r)^2}. \]  
(4.27)

For \( r(u) = \gamma(u, r_0) \), it holds, in view of (4.8) and (4.21), that

\[ r = r_1 + \frac{1}{2} \int_{u_1}^u \tilde{g}(s, r(s)) ds \geq r_1 + \frac{l(x)}{2} (u_1 - u). \]  
(4.28)

Since \( l(x) \in (0, 1) \) for all \( x \in [0, x_0) \), (4.28) implies

\[ 1 + u + r \geq 1 + u + r_1 + \frac{l(x)}{2} (u_1 - u) \geq \frac{l(x)}{2} (1 + u_1 + r_1). \]  
(4.29)

From (4.27) and (4.29), we gain the estimate

\[ \int_0^{u_1} |f(u)| du = \int_0^{u_1} |f(u_1, r_1)| du \leq \frac{(12 + 56K_0)(x^3 + x^5 + x^{2p+1} + x^{2p+3})}{48(1 + u_1 + r_1)^2 l^2(x)}. \]  
(4.30)
Let us now estimate the matrix function \( N \) defined in (4.11) as follows:

\[
|N(u, r)| \leq \left| \frac{1}{2r}(g - \tilde{g}) \right| + \left| \frac{Q^2 g}{4r^3} \right| + \left| \frac{r}{2} Vg \right| + \left| \frac{a}{2} \right|. \tag{4.31}
\]

In view of (4.20), it holds that

\[
\left| \frac{1}{2r}(g - \tilde{g})(u, r) \right| \leq \frac{(3 + 8K_0)(x^2 + x^4 + x^{2p+2})}{48(1 + u)^5}. \tag{4.32}
\]

Since \( 0 < g \leq 1 \), (4.17) yields

\[
\left| \frac{Q^2 g}{4r^3}(u, r) \right| \leq \frac{x^4}{16(1 + u)^4}. \tag{4.33}
\]

From assumption (4.2), since \( p \geq 3/2 \) and \( 0 < g \leq 1 \), we have as in (4.26)

\[
\left| \frac{rg(u, r)}{2} V(\Phi^i \Phi) \right| \leq \frac{K_0 x^{2p+2}}{2(1 + u)^5}. \tag{4.34}
\]

Using (3.4) and (4.17), we estimate \( a/2 \) as

\[
\left| \frac{a(u, r)}{2} \right| \leq \frac{x^2}{2(1 + u)^3}. \tag{4.35}
\]

Inserting (4.32)–(4.35) into (4.31) yields

\[
|N(u, r)| \leq \frac{(30 + 32K_0)(x^2 + x^4 + x^{2p+2})}{48}. \tag{4.36}
\]

From (4.32), we get

\[
\int_0^{u_1} \left| \frac{1}{2r}(g - \tilde{g}) \right|_\gamma \, du \leq \frac{(3 + 8K_0)(x^2 + x^4 + x^{2p+2})}{192}, \tag{4.37}
\]

(4.33) yields

\[
\int_0^{u_1} \left| \frac{Q^2 g}{4r^3} \right|_\gamma \, du \leq \frac{x^4}{48}. \tag{4.38}
\]

From (4.34), we have

\[
\int_0^{u_1} \left| \frac{rg}{2} V \right|_\gamma \, du \leq \frac{K_0 x^{2p+2}}{8}. \tag{4.39}
\]

Using (4.35), we gain

\[
\int_0^{u_1} \left| \frac{a}{2} \right|_\gamma \, du \leq \frac{x^2}{4}. \tag{4.40}
\]
It follows from (4.31), (4.37)–(4.40) that
\[
\left| \int_0^{u_1} [N]_y \, dv \right| \leq K_1 (x^2 + x^4 + x^{2p+2}),
\]
(4.41)
where
\[
K_1 = \frac{55 + 32K_0}{192}.
\]
Using (4.9), it holds that, for \( x \in [0, x_0) \),
\[
1 + r_0 = 1 + r_1 + \frac{1}{2} \int_0^{u_1} \tilde{g}(s, r(s)) \, ds \geq 1 + r_1 + \frac{l(x)}{2} u_1 \geq \frac{l(x)}{2} (1 + r_1 + u_1).
\]
(4.42)
Thus, by the definition of the Banach space \( (\mathcal{X}_0, \| \cdot \|_{\mathcal{X}_0}) \) (see (4.1)),
\[
|W(0, r_0)| \leq \frac{\| W_0 \|_{\mathcal{X}_0}}{(1 + r_0)^2} \leq \frac{4\| W_0 \|_{\mathcal{X}_0}}{(1 + r_1 + u_1)^2 l^2(x)}.
\]
(4.43)
Considering (4.10), (4.30), (4.41) and (4.43), we finally arrive at the following estimate for the solution \( w = \mathcal{K}(W) \) of the initial value problem (4.6)
\[
|w(u_1, r_1)| \leq \frac{K_2 (x^3 + x^5 + x^{2p+1} + x^{2p+3} + \| W_0 \|_{\mathcal{X}_0}) \exp(K_1 (x^2 + x^4 + x^{2p+2})}{(1 + r_1 + u_1)^2 l^2(x)},
\]
(4.44)
where
\[
K_2 = \frac{204 + 56K_0}{48},
\]
(4.44) yields
\[
\sup_{r, u \geq 0} [(1 + u + r)^2 |w(u, r)|] \leq \frac{K_2 (x^3 + x^5 + x^{2p+1} + x^{2p+3} + \| W_0 \|_{\mathcal{X}_0}) \exp(K_1 (x^2 + x^4 + x^{2p+2}))}{l^2(x)}.
\]
(4.45)
We also have to estimate \( \sup_{r, u \geq 0} [(1 + u + r)^3 |w'(u, r)|] \). Set
\[
z(u, r) = w'(u, r), \quad \text{with} \quad z(0, r_0) = W'(0, r_0).
\]
Differentiation of (4.6) w.r.t. \( r \) gives
\[
Dz = N_1 z + B_1 w + B_2 \bar{W} + B_3 \bar{W}',
\]
(4.46)
The spherically symmetric EYMH equations

where the matrix functions $N_1$, $B_1$, $B_2$ and $B_3$ are given by

$$
N_1 = \left[ \frac{\tilde{g}'}{2} + \frac{g - \tilde{g}}{2r} - \frac{Q^2 g}{4r^3} - \frac{rVg'}{2} \right] I_2 + \frac{a}{2} i\sigma_2 = \frac{\tilde{g}'}{2} I_2 + N,
$$

$$
B_1 = \left[ \frac{(g - \tilde{g})'}{2r} - \frac{g - \tilde{g}}{2r^2} - \frac{QQ'g}{2r^3} + \frac{3Q^2 g'}{4r^4} - \frac{Q^2 g'}{4r^3} \right. \\
- \frac{Vg + r(\Phi^i \Phi)^' V'g + rVg'}{2} I_2 + \frac{a'}{2} i\sigma_2,
$$

$$
B_2 = \left[ \frac{QQ'g}{2r^3} - \frac{3Q^2 g}{4r^4} + \frac{Q^2 g'}{2r^2} - \frac{(g - \tilde{g})'}{2r} + \frac{(g - \tilde{g})}{2r^2} \right. \\
+ \frac{Vg + r(\Phi^i \Phi)^' V'g + rVg'}{2} \\
- \frac{gV' + r[V'g' + (\Phi^i \Phi)' V''g]}{2} I_2 + \left( \frac{Q'g + Qg'}{4r} - \frac{Qg}{4r^2} \right) \sigma_1
$$

and

$$
B_3 = \left[ \frac{Q^2 g}{4r^3} - \frac{1}{2r}(g - \tilde{g}) + \frac{r(V - V')g}{2} \right] I_2 + \frac{Qg}{4r} \sigma_1.
$$

Using the characteristics defined earlier, the solution of (4.46) reads

$$
z(u_1, r_1) = \exp \left( \int_{0}^{u_1} [N_1]_\gamma dv \right) W'(0, r_0) + \int_{0}^{u_1} \exp \left( \int_{u}^{u_1} [N_1]_\gamma dv \right) \left[ f_1 \right]_\gamma du,
$$

where

$$
f_1 = B_1 w + B_2 \tilde{W} + B_3 \tilde{W}'.
$$

Using (3.6), we calculate

$$
\frac{1}{2} \tilde{g}' = \frac{g - \tilde{g}}{2r} - \frac{Q^2 g}{4r^3} - \frac{rVg}{2} + \frac{1}{4r^2} \int_{0}^{r} \left( \frac{Q^2 g}{s^3} \right) ds + \frac{1}{2r^2} \int_{0}^{r} s^2 Vg ds.
$$

From (4.16)–(4.19), (4.34) and (4.50), we gain the estimate

$$
\left| \frac{1}{2} \tilde{g}' \right| \leq \frac{(6 + 32K_0)(x^2 + x^4 + x^{2p+2})}{48}.
$$

From (4.36) and (4.51), being $N_1 = (1/2) \tilde{g}' + N$, we gain

$$
|N_1(u, r)| \leq \frac{36 + 64K_0}{48}(x^2 + x^4 + x^{2p+2}).
$$

We now handle $f_1$. From (3.3) and the definition of $g$ in (3.9), we have

$$
|g'| \leq \frac{x^2 r}{4(1 + u)^2(1 + u + r)^4},
$$

(4.51) and (4.53) give

$$
\left| \frac{(g - \tilde{g})'}{2r} \right| \leq \frac{(6 + 16K_0)(x^2 + x^4 + x^{2p+2})}{24(1 + u)^2(1 + u + r)^3}.
$$
(4.20) implies
\[
\left| \frac{g(u, r) - \bar{g}(u, r)}{2r^2} \right| \leq \frac{(3 + 8K_0)(x^2 + x^4 + x^{2p+2})}{48(1 + u)^5(1 + u + r)^3}. \tag{4.55}
\]

In view of (3.5) and the definition of \(\|W\|_X = x\), we have
\[
|Q'(u, r)| \leq \frac{x^2 r}{(1 + u)(1 + u + r)^2}. \tag{4.56}
\]

Because \(0 < g \leq 1\), we deduce from (4.17) and (4.56) that
\[
\left| \frac{QQ'g}{2r^3} \right| \leq \frac{x^4}{4(1 + u)^3(1 + u + r)^5} \tag{4.57}
\]

and
\[
\left| \frac{3Q^2g'}{4r^4} \right| \leq \frac{3x^4}{16(1 + u)^4(1 + u + r)^4}. \tag{4.58}
\]

From (4.17) and (4.53), we obtain
\[
\left| \frac{Q^2g'}{4r^3} \right| \leq \frac{x^6}{64(1 + u)^6(1 + u + r)^6}. \tag{4.59}
\]

It is easy to see from (2.10) and (3.2) that
\[
|((\Phi^t \Phi)'| \leq \frac{2}{r} |\bar{W}(u, r)||W(u, r) - \bar{W}(u, r)|. \tag{4.60}
\]

From (4.2), (4.12), (4.13) and (4.60), we derive the following estimates:
\[
\left| \frac{Vg + r(\Phi^t \Phi)' V'g + rVg'}{2} \right| \leq \frac{K_0x^{2p+2}}{2(1 + u)^5(1 + u + r)^3} + \frac{K_0x^{2p+2}}{2(1 + u)^5(1 + u + r)^3} + \frac{K_0x^{2p+4}}{8(1 + u)^7(1 + u + r)^7}. \tag{4.61}
\]

From (3.4) and (4.17), we have
\[
\left| \frac{a'}{2} \right| = \frac{Qg}{2r^2} \leq \frac{x^2}{4(1 + u)^2(1 + u + r)^2}. \tag{4.62}
\]

Inserting (4.54), (4.55), (4.57)–(4.59), (4.61) and (4.62) into the definition of \(B_1\) in (4.47), we obtain
\[
|B_1| \leq \frac{K_3(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})}{(1 + u)^2(1 + u + r)^2}, \tag{4.63}
\]

where \(K_3 > 0\) is an affine non-decreasing function of \(K_0\). (4.63) yields
\[
|B_1 w(u, r)| \leq \frac{K_3(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})}{(1 + u)^2(1 + u + r)^4} \times \sup_{u, r \geq 0} [(1 + u + r)^2|w(u, r)|]. \tag{4.64}
\]
We now estimate \(|B_2 \bar{W}(u, r)|\). From (4.2), (4.12), (4.13), (4.53) and (4.60), we get
\[
\left| \frac{g V' + r[V' \phi + (\Phi')\phi' V' g]}{2} \right| \leq \frac{9K_0(x^{2p} + x^{2p+2})}{8(1 + u)^3(1 + u + r)^3}.
\] (4.65)

(4.17), (4.53) and (4.56) give
\[
\left| \frac{Q' \phi + Q \phi'}{4r} - \frac{Q \phi}{4r^2} \right| \leq \frac{13(x^2 + x^4)}{32(1 + u)(1 + u + r)^2}.
\] (4.66)

Inserting (4.64), (4.68) and (4.70) into the definition of \(B_2\) in (4.47), we obtain
\[
|B_2| \leq \frac{K_4(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})}{(1 + u)(1 + u + r)^2},
\] (4.67)

where \(K_4 > 0\) is an affine non-decreasing function of \(K_0\). (4.12) and (4.67) give
\[
|B_2 \bar{W}| \leq \frac{K_4(x^3 + x^5 + x^7 + x^{2p+1} + x^{2p+3} + x^{2p+5})}{(1 + u)^2(1 + u + r)^3}.
\] (4.68)

We now handle \(|B_3 \bar{W}'(u, r)|\). From (4.2), (4.17), (4.20) and the definition of \(B_3\) in (4.47), we have
\[
|B_3| \leq \frac{(3 + 14K_0)(x^2 + x^4 + x^{2p} + x^{2p+2})}{12(1 + u)^2(1 + u + r)}.
\] (4.69)

From (4.13) and (4.69), we gain
\[
|B_3 \bar{W}'(u, r)| \leq \frac{(3 + 14K_0)(x^3 + x^5 + x^{2p+1} + x^{2p+3})}{24(1 + u)^3(1 + u + r)^3}.
\] (4.70)

Inserting (4.64), (4.68) and (4.70) into the definition of \(f_1\) in (4.49), and using (4.45), we gain
\[
|f_1(u, r)| \leq \frac{K_2K_3(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})}{(1 + u)^2(1 + u + r)^4l^2(x)} \times (x^3 + x^5 + x^{2p+1} + x^{2p+3} + \|W_0\|_{\mathcal{X}}) \exp(K_1(x^2 + x^4 + x^{2p+2}))
\] 
\[+ \frac{K_5(x^3 + x^5 + x^7 + x^{2p+1} + x^{2p+3} + x^{2p+5})}{(1 + u)^2(1 + u + r)^3}.
\] (4.71)

where \(K_5 > 0\) is an affine non-decreasing function of \(K_0\). Using similar tools as in (4.30), we have
\[
\int_0^{v_1} |[f_1]|\gamma \, du \leq \frac{8K_2K_3(x^3 + x^5 + x^{2p+1} + x^{2p+3} + \|W_0\|_{\mathcal{X}})}{l^5(x)(1 + u_1 + r_1)^3} \times (x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4}) \exp(K_1(x^2 + x^4 + x^{2p+2}))
\] 
\[+ \frac{8K_5(x^3 + x^5 + x^7 + x^{2p+1} + x^{2p+3} + x^{2p+5})}{l^3(x)(1 + u_1 + r_1)^3}.
\] (4.72)
Because $N_1 = (1/2)\tilde{g}' + N$, it follows from (4.41) and (4.51) that
\[
\left| \int_0^{u_1} [N_1]_1 \, dv \right| \leq K_6(x^2 + x^4 + x^{2p+2}),
\]
where $K_6 > 0$ is an affine non-decreasing function of $K_0$. By analogy with (4.43) (see (1.1) for the definition of the Banach space $(X_0, \|\cdot\|_{X_0})$), it holds that
\[
\|W'(0, r_0)\| \leq \frac{8\|W_0\|_{X_0}}{(1 + r_1 + u_1)^3 l^3(x)}.
\]
Considering (4.72)–(4.74), we finally arrive at the following estimate for the solution $z$ of the initial value problem (4.46):
\[
|z(u_1, r_1)| \leq \left[ \frac{8\|W_0\|_{X_0}}{(1 + r_1 + u_1)^3 l^3(x)} + \frac{8K_2K_3(x^3 + x^5 + x^{2p+1} + x^{2p+3} + \|W_0\|_{X_0})}{l^5(x)(1 + u_1 + r_1)^3} \right.
\]
\[
\times \exp(K_7(x^2 + x^4 + x^{2p+2})),
\]
where $K_7 > 0$ is an affine non-decreasing function of $K_0$. From (4.75), we deduce that
\[
\sup_{u, r \geq 0} [(1 + r + u)^3 |z(u, r)|] \leq \left[ \frac{8K_2(x^3 + x^5 + x^{2p+1} + x^{2p+3} + \|W_0\|_{X_0})}{l^3(x)} \times K_3(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4}) \right.
\]
\[
\times \left. \frac{8K_5(x^3 + x^5 + x^7 + x^{2p+1} + x^{2p+3} + x^{2p+5})}{l^3(x)} \exp(K_7(x^2 + x^4 + x^{2p+2})). \right]
\]
Since $l(x) \in (0, 1]$ for all $x \in [0, x_0)$, it follows from (4.45) and (4.76) that
\[
\|w\|_{X'} \leq \frac{K_2(x^3 + x^5 + x^{2p+1} + x^{2p+3} + \|W_0\|_{X_0})}{l^5(x)} \exp(K_9(x^2 + x^4 + x^{2p+2})) \times \frac{K_8(x^3 + x^5 + x^7 + x^{2p+1} + x^{2p+3} + x^{2p+5} + \|W_0\|_{X_0})}{l^5(x)} \exp(K_9(x^2 + x^4 + x^{2p+2})),
\]
where
\[
K_8 = (8 + 8K_4 + 8K_2)K_3 \quad \text{and} \quad K_9 = K_1 + K_7.
\]
The spherically symmetric EYMH equations

Setting

\[ \Omega(x) = \frac{x \ell_5(x) \exp[-K_0(x^2 + x^4 + x^{2p+2})]}{K_2 + K_8(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})} - (x^3 + x^5 + x^{2p+1} + x^{2p+3} + x^{2p+5}), \]

one sees that

\[ \Omega(0) = 0, \quad \Omega'(0) > 0 \quad \text{and} \quad \lim_{x \to +\infty} \Omega(x) = -\infty. \]

Hence, there exists \( x_1 \in (0, x_0) \) such that the function \( \Omega \) is strictly monotonically increasing on \([0, x_1]\). We then choose \( x \in (0, x_1) \) and \( \| W_0 \|_{\mathcal{X}} < x \) to have \( \| w \|_{\mathcal{X}} \leq x \).

This shows that \( \mathcal{K} \) is a mapping from \( B_\rho \) into itself for \( \rho \in (0, x_1) \) and \( \| W_0 \|_{\mathcal{X}} < \rho \).

(ii) The mapping \( \mathcal{K} \) contracts in \( \mathcal{Y} \)

Let \( w_1 \) and \( w_2 \) be two solutions of (4.6) with \( w_1(0, r) = w_2(0, r) = W_0(r), \) \( w_j = \mathcal{K}(W_j), \) \( W_j \in \mathcal{X}, \) \( j = 1, 2. \) We use the following notations: \( g_j = g(W_j), \) \( Q_j = Q(W_j), \)

\( V_j = V(W_j), \) \( z_j = z(W_j), \) \( a_j = a(W_j). \) We assume

\[ \max\{\| W_1 \|_{\mathcal{X}}, \| W_2 \|_{\mathcal{X}}\} \leq x < x_1, \]

and set for convenience

\[ w_1 - w_2 = \vartheta, \quad W_1 - W_2 = \Theta, \quad \| \Theta \|_{\mathcal{Y}} = y \quad \text{and} \quad D_1 = \frac{\partial}{\partial u} - \frac{\tilde{g}_1}{2} \frac{\partial}{\partial r}. \]

Then we have the following nonlinear initial value problem with unknown \( \vartheta \)

\[ D_1 \vartheta = N_2 \vartheta + B_4 \tilde{\Theta} + B_5 (w_2 - \tilde{W}_2) + B_6 \]

and

\[ \vartheta(0, r) = 0, \]

where

\[ N_2 = \left( \frac{g_1 - \tilde{g}_1}{2r} - \frac{Q_1^2 g_1}{4r^3} - \frac{r V_1 g_1}{2} \right) I_2 + \frac{a_1}{2} i \sigma_2, \]

\[ B_4 = \left( \frac{Q_1^2 g_1}{4r^3} - \frac{g_1 - \tilde{g}_1}{2r} + \frac{r g_1 (V_1 - V_1')}{2} \right) I_2 + \frac{Q_1 g_1}{4r} \sigma_1, \]

\[ B_5 = \frac{(g_1 - \tilde{g}_1) - (g_2 - \tilde{g}_2) - (Q_1^2 g_1 - Q_2^2 g_2) - r (V_1 g_1 - V_2 g_2)}{2 \sigma_1 \tilde{W}_2} \]

and

\[ B_6 = -\frac{r (V_1' g_1 - V_2' g_2)}{2} \tilde{W}_2 + \frac{Q_1 g_1 - Q_2 g_2}{4r} \sigma_1 \tilde{W}_2 \]

\[ + \frac{\tilde{g}_1 - \tilde{g}_2}{2} (w_2)' + \frac{a_1 - a_2}{2} i \sigma_2 w_2. \]

Using the characteristic method as in the previous paragraph §4b(ii), we get

\[ \vartheta(u_1, r_1) = \int_0^{u_1} \left\{ \exp \left( \int_u^{u_1} \left[ N_2(u, r) \right]_{\gamma_1} dv \right) \right\} [f_2]_{\gamma_1} \, du, \]

where

\[ f_2 = B_4 \tilde{\Theta} + B_5 (w_2 - \tilde{W}_2) + B_6, \]
and $\gamma_1$ is the characteristic defined by
\[
\frac{dr}{du} = -\frac{1}{2} \tilde{g}_1(u, \gamma_1(u, r)), \quad r(0) = r_0.
\]  
We first estimate $N_2$ as in (4.41) to get
\[
\left|\int_0^{u_1} [N_2] \gamma_1 \, dv\right| \leq K_1(x^2 + x^4 + x^{2p+2}).
\]  
We now concentrate on estimating $f_2$. Similar to (4.27), it is straightforward to have
\[
|B_4 \tilde{Q}| \leq (12 + 56 K_0)(x^2 + x^4 + x^{2p} + x^{2p+2})y.
\]  
Using the mean value theorem and similar calculation as in (4.13), we have
\[
|g_1 - g_2| \leq \frac{xy}{3(1 + u)^2(1 + u + r)^2}
\]  
and
\[
|\tilde{g}_1 - \tilde{g}_2| \leq \frac{xy}{3(1 + u)^3(1 + u + r)}.\]  
In the same way, using the mean value theorem and similar calculation as in (4.12), we have
\[
|Q_1 - Q_2| \leq \frac{xyr^2}{(1 + u)^2(1 + u + r)^2}.
\]  
From (4.17), (4.86)–(4.88), we have
\[
|Q_1^2 g_1 - Q_2^2 g_2| \leq \frac{13(x^3 + x^5)yr^4}{12(1 + u)^4(1 + u + r)^4},
\]  
which implies
\[
\left|\frac{(Q_1^2 g_1 - Q_2^2 g_2)}{4r^3}\right| \leq \frac{13(x^3 + x^5)yr}{48(1 + u)^4(1 + u + r)^4}.
\]  
From (3.6) and the definition of $\tilde{g}$ in (3.9), we have
\[
|\tilde{g}_1 - \tilde{g}_2| \leq |\tilde{g}_1 - \tilde{g}_2| + \frac{1}{2r} \int_0^r \frac{|Q_1^2 g_1 - Q_2^2 g_2|}{s^2} \, ds + \frac{1}{r} \int_0^r s^2 |V_1 g_1 - V_2 g_2| \, ds,
\]  
(4.90) yields
\[
\frac{1}{2r} \int_0^r \frac{|Q_1^2 g_1 - Q_2^2 g_2|}{s^2} \, ds \leq \frac{13(x^3 + x^5)yr^2}{72(1 + u)^5(1 + u + r)^3}.
\]  
Using the mean value theorem and assumption (4.2), we get
\[
|V_1 - V_2| \leq \frac{4K_0 x^{2p+1} y}{(1 + u)^4(1 + u + r)^4}.
\]  

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\[ |V_1 g_1 - V_2 g_2| \leq \frac{(1 + 12K_0)(x^{2p+1} + x^{2p+3})y}{3(1 + u)^4(1 + u + r)^4}. \]  

(4.93)

Thus,

\[ \frac{1}{r} \int_0^r s^2 |V_1 g_1 - V_2 g_2| ds \leq \frac{(1 + 12K_0)(x^{2p+1} + x^{2p+3})r^2 y}{9(1 + u)^3(1 + u + r)^3}. \]  

(4.94)

It follows from (4.87), (4.90), (4.91) and (4.94) that

\[ |\tilde{g}_1 - \tilde{g}_2| \leq \frac{(15 + 32K_0)(x + x^3 + x^5 + x^{2p+1} + x^{2p+3})y}{24(1 + u)^3(1 + u + r)} . \]  

(4.95)

By analogy with (4.16), it holds that

\[ |g_1 - g_2 - (\tilde{g}_1 - \tilde{g}_2)| \leq \frac{2xyr}{(1 + u)^3(1 + u + r)^2}. \]  

(4.96)

Combination of (4.91), (4.94) and (4.96) yields

\[ \left| \frac{(g_1 - \tilde{g}_1) - (g_2 - \tilde{g}_2)}{2r} \right| \leq \frac{(31 + 32K_0)(x + x^3 + x^5 + x^{2p+1} + x^{2p+3})y}{24(1 + u)^3(1 + u + r)^2} . \]  

(4.97)

In view of (4.89), (4.93), (4.97) and the definition of $B_5$ in (4.80), we gain

\[ |B_5| \leq \frac{(83 + 160K_0)(x + x^3 + x^5 + x^{2p+1} + x^{2p+3})y}{24(1 + u)^3(1 + u + r)^2} . \]  

(4.98)

From (4.12) and (4.98), we obtain

\[ |B_5(w_2 - \tilde{W}_2)| \leq \frac{(83 + 160K_0)(x^2 + x^4 + x^6 + x^{2p+2} + x^{2p+4})y}{12(1 + u)^3(1 + u + r)^4} . \]  

(4.99)

We now estimate $B_6$. Using once more the mean value theorem and assumption (4.2), we gain

\[ |V_1' - V_2'| \leq \frac{4K_0 x^{2p-1} y}{(1 + u)^2(1 + u + r)^2} . \]  

(4.100)

(4.2), (4.86) and (4.100) imply

\[ |V_1' g_1 - V_2' g_2| \leq \frac{(1 + 12K_0)(x^{2p-1} + x^{2p+1})y}{3(1 + u)^2(1 + u + r)^2} . \]  

(4.101)

(4.12) and (4.101) yield

\[ \left| r(V_1' g_1 - V_2' g_2) \frac{W_2}{2} \right| \leq \frac{(1 + 12K_0)(x^{2p} + x^{2p+2})y}{6(1 + u)^3(1 + u + r)^2} . \]  

(4.102)

From (4.12), (4.17), (4.86) and (4.88), we derive the following estimate:

\[ \left| \frac{Q_1 g_1 - Q_2 g_2}{4r} \frac{\sigma_1 W_2}{2} \right| \leq \frac{7(x^2 + x^4) y}{24(1 + u)^3(1 + u + r)^2} . \]  

(4.103)
In view of (4.95), we have
\[
\left| \frac{\tilde{g}_1 - \tilde{g}_2}{2}(w_2) \right| \leq \frac{(15 + 32K_0)(x^2 + x^4 + x^6 + x^{2p+2} + x^{2p+4})y}{48(1 + u)^3(1 + u + r)^4}.
\] (4.104)

We use (3.4), (4.86) and (4.88) to gain
\[
|a_1 - a_2| \leq \frac{19(x + x^3)y}{18(1 + u)^3}.
\] (4.105)

Since \(\max\{\|W_1\|_X, \|W_2\|_X\} \leq x < x_1\), (4.105) implies
\[
\left| \frac{a_1 - a_2}{2i\sigma_2w_2} \right| \leq \frac{19(x^2 + x^4)y}{36(1 + u)^3(1 + u + r)^2}.
\] (4.106)

Inserting (4.102)–(4.104) and (4.106) into the definition of \(B_6\) in (4.80), we obtain
\[
|B_6| \leq \frac{(187 + 384K_0)(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})y}{144(1 + u)^3(1 + u + r)^2}.
\] (4.107)

Summing up (4.85), (4.99) and (4.107), and considering the definition of \(f_2\) in (4.82), we arrive at
\[
|f_2| \leq \frac{(1219 + 2472K_0)(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})y}{144(1 + u + r_1)^2l^2(x)}.
\] (4.108)

Thus, using the same tools as in (4.30), we deduce that
\[
\int_0^{u_1} |f_2\gamma_1| \,du \leq \frac{(1219 + 2472K_0)(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})y}{144(1 + u_1 + r_1)^2l^2(x)}.
\] (4.109)

Insertion of (4.84) and (4.109) into (4.81) yields
\[
|\tilde{\vartheta}(u_1, r_1)| \leq \frac{(1219 + 2472K_0)(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})y}{144(1 + u_1 + r_1)^2l^2(x)} \exp(K_1(x^2 + x^4 + x^{2p+2})).
\] (4.110)

Hence
\[
\|\tilde{\vartheta}\|_Y \leq \Xi(x)y,
\] (4.111)

where
\[
\Xi(x) = \frac{(1219 + 2472K_0)(x^2 + x^4 + x^6 + x^{2p} + x^{2p+2} + x^{2p+4})}{144l^2(x)} \exp(K_1(x^2 + x^4 + x^{2p+2})).
\] (4.112)

It is easy to see that the function \(\Xi\) defined in (4.112) is strictly monotonically increasing on \([0, x_1]\) and \(\Xi(0) = 0\). This shows that there exists \(x_2 \in (0, x_1]\) such that \(\Xi(x) < 1/2\) for all \(x \in (0, x_2]\). Thus, in view of (4.111), the mapping \(W \mapsto \mathcal{K}(W)\) contracts in \(\mathcal{Y}\) for \(\|W\|_X \leq x_2\). This concludes the proof of the global existence and uniqueness of classical solution of (3.8).

The decay property (4.4) of the solution is a direct consequence of the definition (4.1) of the Banach spaces \((X, \|\cdot\|_X)\) and \((Y, \|\cdot\|_Y)\).

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Now, from (3.3) and (4.14), one deduces easily that, for each \( r \geq 0 \), \( g \to 1 \) if \( u \to \infty \). In view of (4.20), this implies that, for each \( r \geq 0 \), \( \tilde{g} \to 1 \) if \( u \to \infty \). So, as \( u \to \infty \), the metric given in Bondi coordinates by (2.8) becomes the Minkowski metric. With this, we are done with the proof of theorem 4.1.

**Remark 4.3.**

(i) Theorem 4.1 was stated and proved, under the more restrictive assumption \( p \geq 3 \), by Chae (2003) for the EMH system. Note that the solution obtained here decays more slowly than that of Christodoulou (1986) concerning the spherically symmetric massless Einstein–scalar field system. We found out that this latter fact stems essentially from the estimate (see (4.17))

\[
\left| \frac{Q}{r} \right| \leq \frac{x^2}{2(1 + u)^2(1 + u + r)},
\]

due to the non-vanishing of the local charge \( Q \). Some questions raised in Chae (2003) are thus answered. It would be interesting to find out whether one can use conformal compactification methods of Penrose (1964) to explain slow decaying of the solution via extension by continuity to conformal null infinity.

(ii) It is worth mentioning that assumption (4.2) is not fulfilled by the (classical) self-interaction potential \( V(t) = t^2 \).

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