Nonstatic charged BTZ-like black holes in $N + 1$ dimensions

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We find an exact nonstatic charged BTZ-like solutions, in $(N+1)$-dimensional Einstein gravity in the presence of negative cosmological constant and a nonlinear Maxwell field defined by a power $s$ of the Maxwell invariant, which describes the gravitational collapse of charged null fluid in an anti-de Sitter background. Considering the situation that a charged null fluid injects into the initially an anti-de Sitter spacetime, we show that a black hole form rather than a naked singularity, irrespective of spacetime dimensions, from gravitational collapse in accordance with cosmic censorship conjecture. The structure and locations of the apparent horizons of the black holes are also determined.

It is interesting to see that, in the static limit and when $N = 2$, one can retrieve $2 + 1$ BTZ black hole solutions.

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I. INTRODUCTION

Interest in $2 + 1$ dimensional gravity has been intensify by the discovery of a black hole solution by Banados et al. [1, 2] (referred as BTZ black hole). It enjoys many properties of its counterparts in 3+1 which makes BTZ a suitable model to understand black hole physics in a quite simpler setting. The BTZ black hole is a solution of the $(2+1)$-dimensional Einstein-Maxwell gravity with a negative cosmological constant $\Lambda = -1/\chi^2$. The metric is given by [1, 2]:

$$ds^2 = -N^2(r)dt^2 + \frac{dr^2}{N^2(r)} + r^2d\varphi^2,$$  

with

$$N^2(r) = \frac{r^2}{\chi^2} - \left[ M + 2q^2 \ln\left(\frac{r}{\chi}\right) \right],$$

where $N^2(r)$ is known as the lapse function and $M$ and $q$ are the mass and electric charge of the BTZ black hole, respectively.

The nonlinear electrodynamics models have been proved to be excellent testing grounds in order to elude some problems that occur in the standard Maxwell theory. Indeed, the interest for nonlinear electrodynamics started with the work of Born and Infeld [3]. There has been significant interest about black hole solutions [4–14], including regular black holes [15, 16], with nonlinear electrodynamics source given as an arbitrary power $s$ of the Maxwell invariant, i.e., $(F_{ab}F^{ab})^s$. The nonlinear electrodynamics coupled to general relativity are explored and it has been derived black-hole solutions with interesting asymptotic behaviors [4, 18]. Hendi [14] demonstrated that Einstein-nonlinear Maxwell gravity can generates BTZ-like solutions in $(N + 1)$-dimensional solutions. He showed that the electric field of BTZ-like solutions is the same as $(2 + 1)$-dimensional BTZ black holes, and also their lapse functions are approximately the same.

Over the past decade there has been an increasing interest in the study of black holes, and related objects, in higher (and lower) dimensions, motivated to a large extent by developments in string theory. The aim of the present paper is to obtain nonstatic higher-dimensional analogue of well known BTZ black hole solution in $2 + 1$ dimensions. To be precise, we are interested in $(N+1)$-dimensional exact nonstatic Vaidya-like solution of the Einstein gravity coupled with nonlinear Maxwell source in the presence of null fluid thereby by generalizing those discussions in [14]. The Vaidya geometry permitting the incorporations of the effects of null fluid or null dust offers a more realistic background than static geometries, where all back reaction is ignored. One of the important null fluid solutions include the Robinson - Trautman solution which include, as a special case, the Vaidya solution [13, 19] and in turn include the Schwarzschild vacuum solution. Also, several extension of Vaidya [18, 19] solution, in which the source is a mixture of a perfect fluid and null fluid, have been obtained in later years both in four dimensions [20–22] and higher dimensions [23, 24]. This includes the Bonnor-Vaidya solution [25] for the charge case. The Vaidya solution [18] are widely used as a testing ground for various gravitational scenario and formulation of the cosmic censorship conjecture (CCC) [26, 27]. The CCC proposes that singularities are always hidden within event horizon, and therefore cannot be seen from the rest of spacetime, i.e., no naked singularities. For the ultimate fate of gravitational collapse we could still do no better than the CCC [26, 30]. The CCC, put forward by Penrose 40 years ago, is still one of the most important open questions in general relativity. So far the weak CCC has not been seriously challenged while there exist many counter examples challenging the stronger version of CCC [30].

There are several issues that motivate our analysis: how does the nonlinear Maxwell source change the final fate of collapse? Whether such solutions lead to naked singularities? Do they get covered due to departure from spherically symmetry? Does the nature of the singularity
changes in a more fundamental theory preserving cosmic censorship? As we will see, these collapsing solutions do have several interesting features.

II. FIELD EQUATIONS AND THEIR SOLUTIONS

The \((N+1)\)-dimensional \((N \geq 2)\) action in which gravity is coupled to nonlinear electrodynamics field reads

\[
\mathcal{I} = -\frac{1}{16\pi} \int_M d^{N+1}x \sqrt{-g} (R - 2\Lambda - (\alpha F)^n) + \mathcal{I}_{\text{matter}},
\]

(2)

where \(R\) is scalar curvature, \(\Lambda\) refers to the negative cosmological constant (both \(N + 1\) dimensional), which is in general equal to \(-N(N-1)/2\chi^2\) for asymptotically AdS solutions, in which \(\chi\) is a scale length factor, \(\alpha\) is a constant and \(s\) is power of non-linearity and we choose \(s = N/2\) \([14]\). We consider a null fluid as a matter field, whose action is represented by \(\mathcal{I}_{\text{matter}}\) in Eq. (2).

The field equations obtained by varying the action (2) with respect to the metric \(g_{ab}\) and the gauge field \(A_a\) read respectively

\[
\mathcal{G}_{ab} = R_{ab} - \frac{1}{2} R g_{ab} + \Lambda g_{ab} = T_{ab},
\]

(3)

and

\[
\frac{1}{\sqrt{-g}} \partial_a \left( \sqrt{-g} F^{ab} (\alpha F_{ab F^{ab}})^{N/2 - 1} \right) = 0.
\]

(4)

The Maxwell tensor is \(F_{ab} = \partial_a A_b - \partial_b A_a\), with \(F = F_{ab} F^{ab}\) and \(A_a\) is vector potential. Here the energy-momentum (EMT) tensor is given by

\[
T_{ab} = E_{ab} + T^n_{ab}
\]

(5)

where \(T^n_{ab} = \zeta(v, r) n_a n_b\) with \(\zeta(v, r)\) is the non-zero energy density of null fluid and \(n_a\) is a null vector such that

\[
n_a = \delta_a^0, n_a n^a = 0,
\]

(6)

and the EMT of the nonlinear electrodynamics is defined by

\[
E_{ab} = \alpha (\alpha F)^{N/2 - 1} \left( \frac{1}{2} g_{ab} F - NF_{ab A} F^\lambda \right).
\]

(7)

Now we consider ansatz for the spacetime which expressed in terms of Eddington advanced time coordinate (ingoing coordinate) \(v\) is given by

\[
ds^2 = -e^{\psi(v, r)} dv \left( f(v, r) e^{\psi(v, r)} dv + 2 \epsilon dr \right) + r^2 (d\Omega_{N-1})^2,
\]

(8)

where \((d\Omega_{N-1})^2 = \sum_{i=1}^{N-1} d\theta_i^2\) \(v\) is a null coordinate with \(-\infty < v < \infty\), \(r\) is a radial coordinate with \(0 \leq r < \infty\). Here \(\epsilon = \pm 1\). When \(\epsilon = 1\), the radial coordinate \(r\) increases towards the future along a ray \(v = \text{const}\). When \(\epsilon = -1\), the radial coordinate \(r\) decreases towards the future along a ray \(v = \text{const}\). In what follows, we shall consider \(\epsilon = 1\). We wish to find the general solution of the Einstein equation for the matter field given by Eq. (5) for the metric \([5]\), which contains two arbitrary functions \(\psi(v, r)\) and \(f(v, r)\). It is the field equation \(G^{v}_{v} = 0\) that leads to \(e^{\psi(v, r)} = g(v)\) \([23]\). This could be absorbed by writing \(dv = g(v)dv\). Therefore the entire family of solutions we are searching for is determined by a single function \(f(v, r)\).

Henceforth, in this section, we adopt here a procedure similar to Hendi \([14]\), which we modify here to accommodate our nonstatic case. Without loss of generality, we modify the vector potential, in our case, as

\[
A_a = h(v, r) \delta_a^0.
\]

(9)

In addition, we only consider purely radial electric ansatz do the electromagnetic field which means only non vanishing component of Maxwell tensor are given by

\[
F_{rv} = - F_{vr} = \frac{\partial h}{\partial r}, \quad F = -2 \left( \frac{\partial h}{\partial r} \right)^2,
\]

(10)

which is negative. The sign of the constant \(\alpha\) will ensure the real solutions. The power Maxwell invariant, \((\alpha F)^{N/2}\), may be imaginary for positive \(\alpha\), when \(N/2\) is fractional (for even dimension). Therefore, we set \(\alpha = -1\). For the vector potential \(\psi\), the non-linear electromagnetic field equation \([11]\) leads to

\[
\frac{\partial^2 h}{\partial r^2} + \frac{\partial h}{\partial r} = 0,
\]

(11)

which admits a solution \(h(v, r) = q(v) \ln(r/\chi)\) and hence the only non-vanishing component of Maxwell tensor \(F_{ab}\) takes the form \(F_{rv} = - F_{vr} = q(v)/r\), where \(q(v)\) is arbitrary. It is interesting to note that the expression of the Maxwell tensor \(F_{ab}\) does not depend on the dimension and its value is similar as \((2 + 1)\)-dimensional BTZ solution \([12]\). It is notable that for higher dimensional linear Maxwell field equation, the Maxwell tensor depends on the dimensionality but in our case, the Maxwell tensor is proportional to \(r^{-1}\) in all dimensions. This results are consistent with the corresponding static case \([13]\).

The model considered is obtained from the EMT \([5]\) which is such that the component \(T_{rv}\) must be nonzero and \(T^r_v = T^v_r\) for null energy condition. The nonzero components EMT would read as:

\[
T^r_v = \zeta(v, r), \quad T^v_v = T^r_r = \frac{(N - 1)}{2} g^{N/2} \left( \frac{q(v)}{r} \right)^N, \quad T^\theta_{\theta_1} = T^\theta_{\theta_2} = \ldots = T^\theta_{\theta_{N-2}} = \frac{q^{N/2}}{2} \left( \frac{q(v)}{r} \right)^N.
\]

(12)

The part of the EMT, \(T^{ab}_{\text{non}}\), can be considered as the component of the matter field that moves along the null hypersurfaces \(v = \text{const}\). Type II fluid is the special
case in which the EMT has double null eigenvector. The only observed occurrence of this form EMT is for zero rest-mass field when they represent radiation [31]. The vector field specifies the direction in which the radiation is moving. The physical phenomena which can be modeled by null fluid solutions can be a beam of neutrinos which are assumed for simplicity to be massless. Whereas the EMT of Type I fluid

\[ T^0_0 = \text{diag}[\rho, P_1, P_2, P_3, \ldots, P_N]. \]  

(13)

The EMT of a Type I fluid has only one timelike eigenvector [20, 23, 31]. This is form of the EMT for all observed matter fields with non-zero rest mass and also for the zero rest mass except the Type II fluid discussed above [31].

It is easy to show that the solutions for \( f(v, r) \) is obtained by solving the field equations [3]. The \((v, r)\) component of the equation is integrated to give the general solution as

\[ f(v, r) = \frac{r^2}{\lambda^2} - r^{2-N} \left[ M(v) + 2N/2 q(v) N \ln \left( \frac{r}{\lambda} \right) \right], \]

(14)

where \( M(v) \) and \( q(v) \) are the arbitrary functions which are related to mass and charge parameters, respectively, and which are restricted only by the energy conditions. Thus we have \( N+1 \) dimensional nonstatic solution of Einstein field equations with the metric [9] for a charged null fluid defined by the energy-momentum tensor [5], i.e., we have a kind of \( N+1 \) dimensional charged radiating metric. Whereas when just \( N = 2 \), it reduces to the Vaidya-like black hole solution obtained by Husain [32] in 2 + 1 dimension. Hence, we refer our solution as nonstatic BTZ like solution in 2+1 dimensional Einstein-nonlinear Maxwell gravity representing gravitational collapse of a charged null fluid in an anti-de Sitter spacetime. The metric function \( f(v, r) \), presented here, differ from the linear higher dimensional Bonnor-Vaidya solutions [23, 24]. It is notable that the electric charge term in the linear case is proportional to \( r^{-2(N-2)} \), but in the presented metric function, nonlinear case, this term is logarithmic as in the corresponding static case [14]. As, in the case of static black hole [14], the electric field of our nonstatic BTZ-like solutions is of similar nature as \((2 + 1)\)-dimensional BTZ black holes. From the \((r, v)\) component of field equation and (14), then

\[ \zeta(v, r) = \frac{N - 1}{2N - 1} \left[ \dot{M}(v) + N q(v)^{(N-1)} \dot{q}(v) \ln \left( \frac{r}{\lambda} \right) \right]. \]

(15)

The special case in which \( \dot{M}(v) = 0 \) and \( \dot{q}(v) = 0 \), Eq. (14) after the transformation

\[ dt = dv - \left( \frac{r^2}{\lambda^2} - r^{2-N} \left[ M(v) + 2N/2 q(v) N \ln \left( \frac{r}{\lambda} \right) \right] \right)^{-1} dr, \]

(16)

leads to static black hole found by Hendi [14]. Further in \( 2 + 1 \)-dimensional limit \( (N = 2) \), these solutions reduce to the well known BTZ solution [2].

The vacuum state, namely, what is to be regarded as empty space, is obtained by letting \( M(v) \rightarrow 0 \) and \( q(v) \rightarrow 0 \).

\[ ds^2 = -\frac{r^2}{\lambda^2} dv^2 + 2dvdv + r^2 (d\Omega_{N-1})^2. \]

(17)

When \( M(v) = 1 \) and \( q(v) = 0 \), one gets anti-de Sitter like spacetime

\[ ds^2 = -(1 + \frac{r^2}{\lambda^2}) dv^2 + 2dvdv + r^2 (d\Omega_{N-1})^2. \]

(18)

A. Energy conditions and horizons

The family of solutions discussed here, in general, belongs to Type II fluid defined in [31]. When \( M = q = \) constant, the matter field degenerates to type I fluid [20, 23]. In the rest frame associated with the observer, the energy-density of the matter will be given by

\[ \mu = T^0_0, \quad \rho = -T^i_i = -T^r_r \]

(19)

and the principal pressures are \( P_r = T^i_i \) (no sum convention).

a) The weak energy conditions (WEC): The energy momentum tensor obeys inequality \( T_{ab}u^au^b \geq 0 \) for any time-like vector [31], i.e., \( \zeta \geq 0, \quad \rho \geq 0, \quad P \geq 0 \). Both WEC and SEC, for a Type II fluid, are identical [20, 23].

b) The dominant energy conditions (DEC): For any time-like vector \( w_a, T^{ab}w_b \geq 0 \), and \( T^{ab}w_b \) is non-space-like vector, i.e., \( \zeta \geq 0, \quad \rho \geq P \geq 0 \). It is easy to show that DEC holds as well.

In order to further discuss the physical nature of our solutions, we introduce the kinematical parameters. Following York [33], a null-vector decomposition of the metric is made of the form

\[ g_{ab} = -n_al^a - l_a n_b + \gamma_{ab}, \]

(20)

where,

\[ n_a = \delta_a^v, \quad l_a = \frac{1}{2} f(v, r) \delta_a^r + \delta_a^v, \]

(21)

\[ \gamma_{ab} = r^2 \delta_a^r \delta_b^v + r^2 \sum_{i=1}^{N} \theta_i^2 \delta_a^i \delta_b^i, \]

(22)

\[ l_a n^a = 0 \quad l_a n_a = -1, \quad l^a \gamma_{ab} = 0; \quad \gamma_{ab} n^b = 0, \]

(23)

The expansion of the null rays parameterized by \( v \) is given by

\[ \Theta = \nabla_a v^a - K, \]

(24)
where the $\nabla$ is the covariant derivative and the surface gravity is

$$K = -n^a t^b \nabla_b n_a.$$  \hfill (25)

We now consider the evolution of the apparent horizon (AH) for the metric $[\delta]$. The AH is the outermost marginally trapped surface for the outgoing photons. The AH can be either null or space-like, that is, it can 'move' causally or acausally $[33]$. It gives the equation

$$m(r, v; c_i) = \frac{p}{2} \text{ for } r_{AH}(v; c_i).$$  \hfill (26)

The $c_i$ are any constants that appear in the mass function. The long time $v \rightarrow \infty$ limit of $r_{AH}$ gives the asymptotic radius of the AH as a function of the constants $c_i$. This is a measure of the black hole mass for asymptotically flat or (anti)-de Sitter spacetimes, namely

$$M_{BH} := \lim_{v \rightarrow \infty} r_{AH}(v; c)/2.$$  \hfill (27)

The AHs are defined as surface such that $\Theta = 0$ or $f(v, r) = 0$ $[33]$. Thus

$$r_{I AH} = \chi \exp \left\{ -\frac{1}{N} \text{LambertW}(0, y) - \frac{M(v)}{2^{N/2}q(v)^N} \right\},$$  \hfill (28)

$$r_{O AH} = \chi \exp \left\{ -\frac{1}{N} \text{LambertW}(-1, y) - \frac{M(v)}{2^{N/2}q(v)^N} \right\},$$  \hfill (29)

where

$$y = \left( -\frac{N\chi^{N-2}e^{\frac{NM(v)}{2^{N/2}q(v)^N}}}{2^{N/2}q(v)^N} \right),$$

and the LambertW function satisfies

$$\text{LambertW}(x) \exp [\text{LambertW}(x)] = x.$$  \hfill (30)

The pure charged case ($M(v) = 0$) is also important, since then

$$f(v, r) = -\frac{r^2}{\chi^2} + r^{(2-N)2^{N/2}}q(v)^N \ln \left( \frac{r}{\chi} \right)$$  \hfill (31)

and we have horizons without mass

$$r_{I AH} = \chi \exp \left\{ -\frac{1}{N} \text{LambertW}(0, -\frac{N-N}{q(v)N/2}) \right\},$$

$$r_{O AH} = \chi \exp \left\{ -\frac{1}{N} \text{LambertW}(-1, -\frac{N-N}{q(v)N/2}) \right\}.$$

The uncharged case ($q(v) \rightarrow 0$) is interesting, since then and apparent horizon is $r_{AH} = \exp\left(\frac{\ln(M(v)\chi^2)}{N}\right)$. It is clear that presence of the nonlinear Maxwell source produces a change in the location of the AHs. Such a change could have a significant effect in the dynamical evolution of the black hole horizon. The timelike limit surface (TLS) for a black hole, with a small luminosity, is locus where $g_{vv} = 0$ $[33]$. Here one sees that $\Theta = 0$, implies $f(v, r_{AH}) = 0$ or $g_{vv}(r = r_{AH}) = 0$ implies that $r = r_{AH}$ is TLS and thus in our case AH and TLS coincide.

### III. CAUSAL STRUCTURE OF SINGULARITIES

The easiest way to detect a singularity in a space-time is to observe the divergence of some invariants of the Riemann tensor. The Kretschmann scalar ($K = R_{abcd}R^{abcd}$, $R_{abcd}$ is the $N + 1$ Riemann tensor) for the metric $[\delta]$ reduces to

$$K = \left( \frac{\partial f}{\partial r^2} \right)^2 + \frac{2(N-1)}{r^4} \left( \frac{\partial f}{\partial r} \right)^2 + \frac{2(N-2)(N-1)}{r^4} \left( f^2 \right),$$  \hfill (33)

which diverges and so is energy density $[15]$ at $r = 0$ for $f \neq 0$ and $N \geq 2$ indicating presence of scalar polynomial singularity. The physical situation is that of a radial inflow of charged null fluid in the region of the anti-de Sitter universe. The first shell arrives at $r = 0$ at time $v = 0$ and the final at $v = T$. A central singularity of growing mass developed at $r = 0$. For $v < 0$ we have $M(v) = q(v) = 0$, i.e., $N + 1$ dimensional vacuum metric $[17]$, and for $v > T$, $M(v) = q(v) = 0$, $M(v)$ and $q(v)$ are positive definite. The metric for $v = 0$ to $v = T$ is $N + 1$ dimensional nonstatic BTZ-like radiating solution derived above, and for $v > T$ we have the $N + 1$ dimensional static BTZ-like solution $[7]$

$$ds^2 = \left( \frac{r^2}{\chi^2} - r^{2-N} \left[ M + 2^{N/2}q(v)\ln \left( \frac{r}{\chi} \right) \right] \right) \text{d}v^2 + 2dvdr$$

$$+ r^2 (d\Omega_{N-1}^2).$$  \hfill (34)

Radial ($\theta$ and $\phi = \text{const}$.) null geodesics of the metric (11) must satisfy the null condition

$$\frac{dv}{dr} = \frac{2}{\left( \frac{r^2}{\chi^2} - r^{2-N} \left[ M + 2^{N/2}q(v)\ln \left( \frac{r}{\chi} \right) \right] \right)}.$$  \hfill (35)

Clearly, the above differential equation has a singularity at $r = 0$, $v = 0$. The nature (a naked singularity or a black hole) of the collapsing solutions can be characterized by the existence of radial null geodesics coming out from the singularity. The nature of the singularity can be analyzed by techniques in $[27]$. To proceed further, we choose

$$M(v) = \lambda v^{(N-2)} (\lambda > 0) \text{ and } 2^{N/2}q(v)^N = \mu^2 v^{(N-2)} (\mu > 0),$$

for $0 \leq v \leq T$ $[26, 34]$. Let $Y = v/r$ be the tangent to a possible outgoing geodesic from the singularity.
In order to determine the nature of the limiting value of $Y$ at $r = 0$, $v = 0$ on a singular geodesic, we let $Y_0 = \lim_{r \to 0} v \to 0 Y = \lim_{r \to 0} v \to 0 \frac{\nu}{r}$.

Using \( \text{E} \) and L'Hôpital's rule we get

\[
Y_0 = \lim_{r \to 0} \frac{v}{r} = \lim_{r \to 0} \frac{\nu}{r} = \lim_{r \to 0} \frac{dv}{dr} = \lim_{r \to 0} \frac{2}{v} N - 2 \left( Y N - 2 + \mu^2 Y N - 2 \ln \left( \frac{x}{\lambda} \right) \right).
\]

This is the equation which would ultimately decide the end state of collapse: a black hole or a naked singularity. Thus by analyzing this algebraic equation, the nature of the singularity can be determined. The central shell focusing singularity would atleast be locally naked (for brevity we have addressed it as naked throughout this paper), if Eq. \( \text{E} \) admits one or more positive real roots \([26]\). The values of the roots give the tangents of the escaping geodesics near the singularity. When there are no positive real roots to Eq. \( \text{E} \), there are no outgoing future directed null geodesics emanating from the singularity. Thus, the occurrence of positive roots would imply the violation of the strong CCC, though not necessarily of the weak form. Hence in the absence of positive real roots, the collapse will always lead to a black hole. Clearly, Eq. \( \text{E} \) do not admit positive roots \([26]\) and no radial future null geodesics terminate at the singularity. Thus referring to the above discussion, the collapse proceed to form a black hole.

\section{A. CLOSING REMARKS}

In this paper we have demonstrated a construction \( N + 1 \) dimension charged BTZ-like nonstatic black hole solutions, namely \( N + 1 \) dimensional Vaidya-form of BTZ black holes. We have obtained the general black hole solutions for charged null fluid for the metric \textit{ansatz} \([8]\) in the Eddington advanced time coordinate of \( N + 1 \) dimensional gravity coupled with a nonlinear electrodynamics theory as a power of the Maxwell invariant. This yields, in \( 2 + 1 \) dimension and in static limit, the BTZ solutions \([2]\). We have used the solution to discuss the consequence of nonlinear electrodynamics on the structure and location of the horizons these radiating black holes. The AHs are obtained exactly by method developed by York \([33]\) and it is also shown that the AHs of these black holes coincides with TLS as it should be.

The gravitational collapse of spherical matter in the form of null fluid described by Vaidya metric \([18]\) is well studied for investigating CCC. It clear that the central shell focusing singularity can be naked or covered depending upon on the choice of initial data. The scenarios considered so far (see \([26]\) for details) are spherically symmetric and asymptotically flat. We may then ask if the occurrence of a naked singularity in these cases is an artefact of the special symmetry. In the presence of negative cosmological term and departure from spherical symmetry one can expect the occurrence of major changes. When a negative cosmological constant is introduced, the spacetime will become asymptotically anti-de Sitter spacetime. The present case is an example of non spherical symmetry as well as non asymptotic flatness, and we have shown both uncharged and charged collapse of null fluid does not yield naked singularities, but always black holes. This shows that non-spherical collapse of null fluid (charged or uncharged) discussed here in a negative cosmological background does nor violate CCC and the result is independent of spacetime dimension. We therefore conclude that CCC is actually respected for our nonstatic BTZ-like solution in all dimensions.

In conclusion, we have derived a new class of \( N + 1 \) dimensional charged null fluid collapse solutions of the Einstein equations with a negative cosmological constant. As a special case these solutions can be reduced to the BTZ black hole. The study of geometric properties, causal structures and thermodynamics of these black hole solutions will be subject of the future project. As final remark it would be interesting to explore the extensions of the solutions presented here in more general context, e.g., to see how the results get modified with the inclusion of Gauss-Bonnet combination of quadratic curvature terms and, in general, for the Lovelock polynomial in the action \([2]\), which is left for future investigation \([33]\).

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