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Black holes in an expanding universe and supersymmetry

Dietmar Klemm, Masato Nozawa *

Dipartimento di Fisica, Università di Milano, and INFN, Sezione di Milano, Via G. Celoria 16, 20133 Milano, Italy

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ABSTRACT

This paper analyzes the supersymmetric solutions to five and six-dimensional minimal (un)gauged supergravities for which the bilinear Killing vector constructed from the Killing spinor is null. We focus on the spacetimes which admit an additional SO(1, 1) boost symmetry. Upon the toroidal dimensional reduction along the Killing vector corresponding to the boost, we show that the solution in the ungauged case describes a charged, nonextremal black hole in a Friedmann–Lemaître–Robertson–Walker (FLRW) universe with an expansion driven by a massless scalar field. For the gauged case, the solution corresponds to a charged, nonextremal black hole embedded conformally into a Kantowski–Sachs universe. It turns out that these dimensional reductions break supersymmetry since the bilinear Killing vector and the Killing vector corresponding to the boost fail to commute. This represents a new mechanism of supersymmetry breaking that has not been considered in the literature before.

1. Introduction

Over the last two decades, many developments of superstring theory have been triggered by supersymmetric solutions in supergravities. In particular, supersymmetric black holes played a key role for the first successful account for the microscopic origin of the Bekenstein–Hawking entropy [1]. Recently a systematic classification of supersymmetric solutions has been developed and proved useful for obtaining supersymmetric black objects with various topologies (see e.g. [2–13] for an incomplete list). The supersymmetric solutions are divided into two categories, according to the causal character of the vector field constructed from the Killing spinor, i.e., timelike and null classes. Typically, the timelike class of solutions contains black holes, whereas the null family contains propagating waves. The timelike class of metrics in ungauged supergravities is specified by a set of harmonic/Poisson-type functions on a $(d-1)$-dimensional manifold with reduced holonomy over which the metric is fibered. It therefore follows that supersymmetric black holes belonging to the timelike class are time-independent with degenerate horizons and allow for a superposition principle, as inferred from the Majumdar–Papapetrou solution. This represents a situation in which gravitational and electromagnetic fields are in mechanical equilibrium.

More than twenty years ago, Kastor and Traschen discovered an interesting generalization of the Majumdar–Papapetrou solution in the Einstein–Maxwell-$\Lambda(>0)$ system [14]. The Kastor–Traschen solution is characterized by a harmonic function on $\mathbb{R}^3$ with an additional time-dependence and asymptotically tends to the de Sitter universe. When the harmonic function has a single monopole source at the center of $\mathbb{R}^3$, the metric describes a black hole with a bifurcate Killing horizon in the de Sitter universe, i.e., the lukewarm limit of the Schwarzschild–de Sitter black hole [15]. The superposition property of the Kastor–Traschen solution is reminiscent of supersymmetric solutions in the timelike class, although a positive cosmological constant is not compatible with supersymmetry. Nevertheless, the Kastor–Traschen solution admits a spinor obeying 1st-order differential equations in "fake" supergravity, in which the gauge coupling constant in gauged supergravity is analytically continued [16,17]. The superposition property further allows to investigate analytically black hole collisions in a (contracting) universe and to test the validity of the cosmic censorship conjecture [18].

Later on, Ref. [19] obtained a time-dependent and spatially inhomogeneous solution from the time-dependent intersecting M2/M2/M5/M5 branes, which reduces to $AdS_2 \times S^2$ for $r \to 0$, and approaches for $r \to \infty$ to the FLRW cosmology with the scale factor obeying $a(\tau) \propto \tau^{1/3}$. Maeda and one of the present authors verified that this metric indeed describes a black hole in the

* Corresponding author.
E-mail addresses: dietmar.klemm@mi.infn.it (D. Klemm), masato.nozawa@mi.infn.it (M. Nozawa).

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FLRW universe with regular horizons [20]. The solution was further generalized to the case with a Liouville-type scalar potential, for which the metric asymptotically tends to an FLRW universe with arbitrary power-law expansion [21,22]. These solutions are very similar to the Kastor–Traschen solution since they are specified by some set of harmonic functions on a base space. Interestingly, the event horizon is generated by an asymptotic Killing vector and realizes the isolated horizon [23], when each harmonic has a point source at the origin. Hence, the area of the horizon fails to grow even though the outside region of the black hole is highly dynamical. Moreover, it was shown that these solutions are pseudo-supersymmetric in “fake” supergravity [24]. Using the general classification scheme of [25], further extensions to the case with a sum of exponential scalar potentials and to the case including rotation were analyzed in Refs. [26,27].

The cosmic expansion of the solution in Ref. [19] is driven by a massless scalar field corresponding to a “flat gauging” in the context of gauged supergravity. It might therefore be possible to embed these solutions into higher-dimensional supersymmetric spacetimes by the Kaluza–Klein mechanism, rather than embedding them into fake supergravity. As we commented, a naive Kaluza–Klein reduction does not work, since supersymmetric black holes are time-independent and extremal, whereas the solution in [19] is time-evolving and non-extremal. To fill this gap is one of the main aims of the present article.

We exhibit a class of supersymmetric solutions which can be identified as a black hole in an expanding universe upon dimensional reduction. Interestingly, the black hole is time-dependent and admits nondegenerate horizons, both of these properties counter to those for supersymmetric black holes in the timelike class. This is possible because our supersymmetric solutions belong to the null family. We discuss how an additional SO(1, 1) scaling property gives rise to a Killing vector for the dimensional reduction and how this Kaluza–Klein reduction breaks supersymmetry. This susy breaking mechanism is new, and may have applications in other contexts as well.

The remainder of our paper is organized as follows. In the next section, we show that the five-dimensional null BPS family in minimal (un)gauged supergravity admits solutions describing (after a KK reduction) a black hole in equilibrium in an expanding universe. In section 3, we show how to obtain five-dimensional dynamical black holes from a supersymmetric solution in six-dimensional minimal ungauged supergravity. Section 4 contains our conclusions. We employ the mostly plus metric signature throughout the article.

2. Black hole from five dimensions

2.1. Ungauged case

The bosonic Lagrangian of five-dimensional ungauged minimal supergravity is given by [3]

$$\mathcal{L}_5^{(0)} = R \star 1 - 2F\wedge\star F - \frac{8}{3\sqrt{3}} F\wedge F\wedge A\,, \tag{2.1}$$

where $F = \mathrm{d}A$ is a Maxwell field. In terms of a Dirac spinor ϵ, the Killing spinor equation reads

$$\hat\nabla_\mu \epsilon \equiv \left[\nabla_\mu + \frac{i}{4\sqrt{3}}\left(\gamma_\mu{}^{\nu\rho} - 4\delta_\mu{}^\nu\gamma^\rho\right)F_{\nu\rho}\right]\epsilon = 0\,. \tag{2.2}$$

Let us consider the case in which $V^\mu \equiv i\bar\epsilon\gamma^\mu\epsilon$ is a null vector. In the coordinate system $V = \partial/\partial v$, the metric and the gauge field are v-independent and the general supersymmetric solution in the null family is given by [3]

$$\mathrm{d}s^2 = -2e^+e^- + e^ie^i\,, \qquad A = -\frac{\sqrt{3}}{2}\tilde A_i \mathrm{d}x^i\,, \tag{2.3}$$

where $i, j \ldots = 1, 2, 3$ and the orthonormal frame is given by

$$e^+ = H^{-1}\mathrm{d}u\,, \quad e^- = \mathrm{d}v + \frac{\mathcal{F}}{2}\mathrm{d}u\,, \quad e^i = H(\mathrm{d}x^i + a^i\mathrm{d}u)\,.$$

In three-dimensional vector notation, the supersymmetric solutions are determined by the system

$$\nabla\times\tilde{\mathbf{A}} = \nabla H\,, \qquad \partial_u\tilde{\mathbf{A}} = \frac{1}{3}H^{-2}\nabla\times\left(H^3\mathbf{a}\right)\,,$$
$$\nabla^2\mathcal{F} = 2H^2D_uW_{ii} + 2HW_{(ij)}W_{(ij)} + \frac{2}{3}HW_{[ij]}W_{[ij]}\,, \tag{2.4}$$

where $D_u \equiv \partial_u - \mathbf{a}\cdot\nabla$ and $W_{ij} \equiv D_uH\delta_{ij} - H\partial_j a_i$. The integrability condition of (2.4) leads to $\nabla^2 H = 0$. The solution to the Killing spinor equation (2.2) is given by the constant spinor under the projection $\gamma^+\epsilon = 0$, viz, the solution preserves half of the supersymmetries.

Let us focus here on the following class of supersymmetric solutions

$$\mathbf{a} = 0\,, \qquad H = H(\mathbf{x})\,, \qquad \mathcal{F} = -\frac{4}{(hu)^2}U(\mathbf{x})\,. \tag{2.5}$$

With these restrictions, the metric is invariant under the SO(1, 1) boost action $u \to \lambda u$, $v \to v/\lambda$ [28]. Namely there exists an additional Killing vector $\xi = u\partial/\partial u - v\partial/\partial v$ corresponding to the scaling. By the following coordinate transformation $(u, v) \to (t, w)$:

$$u = \frac{2}{h}e^{-hw/2}\,, \qquad v = te^{hw/2}\,, \tag{2.6}$$

where h is a constant, the scaling Killing vector is transformed into a coordinate vector, $\xi = -(2/h)\partial/\partial w$. It therefore follows that the metric (2.3) is independent of w and reads

$$\mathrm{d}s^2 = H^{-1}\mathrm{d}w[2\mathrm{d}t + (ht+U)\mathrm{d}w] + H^2\mathrm{d}\mathbf{x}^2\,, \tag{2.7}$$

where H and U obey Laplace's equations $\nabla^2 H = \nabla^2 U = 0$ on $\mathbb{R}^3$. One can then reduce the system down to four dimensions by the Kaluza–Klein ansatz

$$\mathrm{d}s^2 = e^{-2\phi/\sqrt{3}}(\mathrm{d}w + 2A^{(1)})^2 + e^{\phi/\sqrt{3}}g_{\mu\nu}\mathrm{d}x^\mu\mathrm{d}x^\nu\,, \tag{2.8}$$

where

$$\phi = \frac{\sqrt{3}}{2}\ln\left(\frac{H}{ht+U}\right)\,, \qquad A^{(1)} = \frac{\mathrm{d}t}{2(ht+U)}\,, \tag{2.9}$$

and the 4-dimensional metric $\mathrm{d}s_4^2 = g_{\mu\nu}\mathrm{d}x^\mu\mathrm{d}x^\nu$ reads

$$\mathrm{d}s_4^2 = -\Xi_4^{-1}\mathrm{d}t^2 + \Xi_4\mathrm{d}\mathbf{x}^2\,, \tag{2.10}$$

with $\Xi_4 \equiv [(ht+U)H^3]^{1/2}$. This recovers the solution obtained by the compactification of dynamically intersecting branes (with three equal charges) [19] and solves the four-dimensional field equations derived from the Lagrangian

$$\mathcal{L}_4^{(0)} = R - \frac{1}{2}(\nabla\phi)^2 - e^{-\sqrt{3}\phi}F^{(1)}_{\mu\nu}F^{(1)\mu\nu} - e^{-\phi/\sqrt{3}}F^{(2)}_{\mu\nu}F^{(2)\mu\nu}\,, \tag{2.11}$$

where $F^{(1,2)} = \mathrm{d}A^{(1,2)}$ and $A^{(2)} = -\frac{\sqrt{3}}{2}\tilde A_i\mathrm{d}x^i$ descends from the five-dimensional gauge potential (2.3).

Working in spherical coordinates $\mathrm{d}\mathbf{x}^2 = \mathrm{d}r^2 + r^2(\mathrm{d}\theta^2 + \sin^2\theta\mathrm{d}\phi^2)$, let us consider the case in which only the monopole sources are nonvanishing as $H = 1 + Q/r$ and $U = Q/r$. Asymptotically for $r \to \infty$, the metric (2.10) then tends to an expanding