Asymptotically de Sitter and anti-de Sitter Black Holes with Confining Electric Potential

Eduardo Guendelman\textsuperscript{a,*}, Alexander Kaganovich\textsuperscript{a}, Emil Nissimov\textsuperscript{b}, Svetlana Pacheva\textsuperscript{b}

\textsuperscript{a}Department of Physics, Ben-Gurion University of the Negev, P.O.Box 653, IL-84105 Beer-Sheva, Israel
\textsuperscript{b}Institute for Nuclear Research and Nuclear Energy, Bulgarian Academy of Sciences, Boul. Tsarigradsko Chausee 72, BG-1784 Sofia, Bulgaria

Abstract

We study gravity interacting with a special kind of QCD-inspired nonlinear gauge field system which earlier was shown to yield confinement-type effective potential (the “Cornell potential”) between charged fermions (“quarks”) in flat space-time. We find new static spherically symmetric solutions generalizing the usual Reissner-Nordström-de-Sitter and Reissner-Nordström-anti-de-Sitter black holes with the following additional properties: (i) appearance of a constant radial electric field (in addition to the Coulomb one); (ii) novel mechanism of dynamical generation of cosmological constant through the non-Maxwell gauge field dynamics; (iii) appearance of confining-type effective potential in charged test particle dynamics in the above black hole backgrounds.

Keywords: black holes in (anti) de Sitter spaces, dynamically generated cosmological constant, QCD-inspired confining potential

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*Corresponding author – tel. +972-8-647-2508, fax +972-8-647-2904.
Email addresses: guendel@bgu.ac.il (Eduardo Guendelman), alexk@bgu.ac.il (Alexander Kaganovich), nissimov@inrne.bas.bg (Emil Nissimov), svetlana@inrne.bas.bg (Svetlana Pacheva)
URL: http://eduardo.hostoi.com (Eduardo Guendelman), http://profiler.bgu.ac.il/frontoffice/ShowUser.aspx?id=1249 (Alexander Kaganovich), http://theo.inrne.bas.bg/~nissimov/ (Emil Nissimov), http://theo.inrne.bas.bg/~svetlana/ (Svetlana Pacheva)
1. Introduction

It has been shown by 't Hooft [1] that in any effective quantum theory, which is able to describe linear confinement phenomena, the energy density of electrostatic field configurations should be a linear function of the electric displacement field in the infrared region due to appropriate infrared counterterms. The simplest way to achieve this in Minkowski space-time is by considering a square root of the field strength squared, in addition to the standard Maxwell term, leading to a very peculiar non-Maxwell nonlinear effective gauge field model [2]:

\[ S = \int d^4x L (F^2), \quad L (F^2) = -\frac{1}{4} F^2 - f \frac{\sqrt{-F^2}}{2}, \]

with \( f \) being a positive coupling constant. It has been shown in first three refs. [2] that the square root of the Maxwell term naturally arises as a result of spontaneous breakdown of scale symmetry of the original scale-invariant Maxwell theory with \( f \) appearing as an integration constant responsible for the latter spontaneous breakdown. The model (1) produces a confining effective potential \( V(r) = -\frac{\alpha}{r} + \beta r \) (Coulomb plus linear one) which is of the form of the well-known “Cornell” potential [3] in quantum chromodynamics (QCD). For static field configurations the model (1) yields the following electric displacement field \( \vec{D} = \vec{E} - \frac{f}{\sqrt{2} |\vec{E}|} \). The pertinent energy density turns out to be (there is no contribution from the square-root term in (1))

\[ \frac{1}{2} E^2 = \frac{1}{2} |\vec{D}|^2 + \frac{f}{\sqrt{2} |\vec{E}|} \frac{\vec{E} \cdot \vec{D}}{|\vec{E}|} + \frac{f^2}{4}, \]

so that it indeed contains a term linear w.r.t. \( |\vec{D}|\).

It is crucial to stress that the Lagrangian \( L(F^2) \) (1) contains both the usual Maxwell term as well as a non-analytic function of \( F^2 \) and thus it is a non-standard form of nonlinear electrodynamics. In this way it is significantly different from the original “square root” Lagrangian \( -\frac{f}{2} \sqrt{F^2} \) first proposed by Nielsen and Olesen [4] to describe string dynamics. Moreover, it is important that the square root term in (1) is in the “electrically” dominated form \( (\sqrt{-F^2}) \) as opposed to the “magnetically” dominated Nielsen-Olesen form \( (\sqrt{F^2}) \).

Let us remark that one could start with the non-Abelian version of the action (1). Since we will be interested in static spherically symmetric solutions,
the non-Abelian theory effectively reduces to an Abelian one as pointed out in the first ref. [2].

Our main goal in the present note is to study possible new effects by coupling the confining potential generating nonlinear gauge field system (1) to gravity. We find:

(i) appearance of a constant radial electric field (in addition to the Coulomb one) in charged black holes within Reissner-Nordström-de-Sitter and/or Reissner-Nordström-anti-de-Sitter space-times as well as in electrically neutral black holes with Schwarzschild-de-Sitter and/or Schwarzschild-anti-de-Sitter geometry;

(ii) novel mechanism of dynamical generation of cosmological constant through the non-Maxwell gauge field dynamics of the nonlinear action $L(F^2)$ (1);

(iii) appearance of confining-type effective potential in charged test particle dynamics in the above black hole backgrounds.

2. Lagrangian Formulation. Spherically Symmetric Solutions

We will consider the simplest coupling of the nonlinear gauge field system (1) to gravity described by the action (we use units with Newton constant $G_N = 1$):

$$S = \int d^4x \sqrt{-g} \left[ \frac{R(g)}{16\pi} - \frac{1}{4} F^2 - \frac{f}{2} \sqrt{-F^2} \right], \quad F^2 \equiv F_{\kappa\lambda} F^{\mu\nu} g^{\kappa\mu} g^{\lambda\nu},$$

(2)

where $R(g)$ is the scalar curvature of the space-time metric $g_{\mu\nu}$ and $g \equiv \det |g_{\mu\nu}|$. It is important to stress that for the time being we will not introduce any bare cosmological constant term.

The energy-momentum tensor $T^{(F)}_{\mu\nu}$ of the nonlinear gauge field, which appears in the pertinent equations of motion:

$$R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R = 8\pi T^{(F)}_{\mu\nu},$$

(3)

$$\partial_{\nu} \left( \sqrt{-g} \left( 1 - \frac{f}{\sqrt{-F^2}} \right) F_{\kappa\lambda} g^{\mu\kappa} g^{\nu\lambda} \right) = 0,$$

(4)

is explicitly given by:

$$T^{(F)}_{\mu\nu} = \left( 1 - \frac{f}{\sqrt{-F^2}} \right) F_{\mu\kappa} F_{\nu\lambda} g^{\kappa\lambda} - \frac{1}{4} \left( F^2 + 2f \sqrt{-F^2} \right) g_{\mu\nu}.$$

(5)
We will look for static spherically symmetric solutions of the system (3)–(5):

\[
ds^2 = -A(r)dt^2 + \frac{dr^2}{A(r)} + r^2(d\theta^2 + \sin^2 \theta d\varphi^2),
\]

(6)

\[
F_{\mu\nu} = 0 \text{ for } (\mu, \nu) \neq (0, r) , \quad F_{0r} = F_{0r}(r).
\]

(7)

In this case the gauge field equations of motion (4) become:

\[
\partial_r \left( r^2 \left( F_{0r} - \frac{\varepsilon_F f}{\sqrt{2}} \right) \right) = 0 , \quad \varepsilon_F \equiv \text{sign}(F_{0r}),
\]

(8)

whose solution reads:

\[
F_{0r} = \frac{\varepsilon_F f}{\sqrt{2}} + \frac{Q}{\sqrt{4\pi r^2}} , \quad \varepsilon_F = \text{sign}(Q).
\]

(9)

Again, as in the flat space-time case (1), the electric field contain a radial constant piece \(\varepsilon_f f/\sqrt{2}\) alongside with the Coulomb term.

Further, it has been shown in ref. [5] that for static spherically symmetric metrics (6) with the associated energy-momentum tensor obeying the condition \(T_{00} = T_{rr}\), which is fulfilled in the present case (7), it is sufficient to solve Einstein equations:

\[
R_{00} = 8\pi \left( T_{00} - \frac{1}{2} T^{\lambda}_{\lambda} \right) \quad \text{where} \quad R_{00} = -\frac{1}{2r^2} \partial_r (r^2 \partial_r A),
\]

\[
R_{\theta\theta} = -8\pi T_{00} \quad \text{where} \quad R_{\theta\theta} = -\frac{1}{r^2} (A - 1) - \frac{1}{r} \partial_r A.
\]

(10)

In the case under consideration the r.h.s. of the first Einstein equation (10) with the energy-momentum tensor (5) becomes:

\[
8\pi \left( T^{(F)}_{00} - \frac{1}{2} T^{(F)}_{\lambda\lambda} \right) = -4\pi \left( \frac{Q^2}{4\pi r^4} - \frac{1}{2} f^2 \right)
\]

(11)

taking into account (7) and (9). Interestingly enough, there are no cross terms in (11) between the Coulomb and constant electric parts.

The solution of (10) with (11) yields:

\[
A(r) = 1 - \sqrt{8\pi |Q| f - \frac{2m}{r} + \frac{Q^2}{r^2} - \frac{2\pi f^2}{3} r^2}.
\]

(12)

In other words the solution given by (6), (12) and (9) describes a black hole with:
• Reissner-Nordström-de-Sitter space-time geometry (12);
• additional global constant radial electric field in (9) apart from the usual Coulomb one;
• dynamically generated effective cosmological constant in (12) (let us recall that there was no bare cosmological constant in (2)):
\[ \Lambda_{\text{eff}} = 2\pi f^2 . \] (13)

In particular, when \( Q = 0 \) we obtain electrically neutral black hole with Schwarzschild-de-Sitter geometry:
\[ A(r) = 1 - \frac{2m}{r} - \frac{2\pi f^2}{3} r^2 . \] (14)
where the cosmological constant (13) is dynamically generated, and with additional global constant radial electric field:
\[ F_{0r} = \varepsilon f / \sqrt{2} . \] (15)

3. Bare Negative Cosmological Constant versus Induced Cosmological Constant

Let us now introduce in (2) from the very beginning a negative bare cosmological constant \( \Lambda = -|\Lambda| \):
\[ S = \int d^4x \sqrt{-g} \left[ \frac{1}{16\pi} (R(g) - 2\Lambda) - \frac{1}{4} F^2 - f \sqrt{-F^2} \right] . \] (16)
Then the corresponding static spherically symmetric solution is given by (9) and (6) with:
\[ A(r) = 1 - \sqrt{8\pi |Q| f} - \frac{2m}{r} + \frac{Q^2}{r^2} + \frac{1}{3} (|\Lambda| - 2\pi f^2) r^2 . \] (17)
Thus, we find also black hole solution with Reissner-Nordström-anti-de-Sitter geometry (17) and with additional global constant electric field (9) provided the full effective cosmological constant (bare one plus dynamically induced one) satisfies:
\[ \Lambda_{\text{eff}} = -|\Lambda| + 2\pi f^2 < 0 \quad , \quad \text{i.e.} \quad |\Lambda| > 2\pi f^2 . \] (18)
On the other hand, if $|\Lambda| < 2\pi f^2$, i.e. $\Lambda_{\text{eff}} = 2\pi f^2 - |\Lambda| > 0$, the solution (17) describes asymptotically de Sitter black hole in spite of the presence of negative bare cosmological constant $\Lambda$. In the special case $|\Lambda| = 2\pi f^2$ the dynamically induced cosmological constant completely cancels the effect of the negative bare cosmological constant and the resulting solution describes an asymptotically flat Reissner-Nordström-like black hole ($A(r) = 1 - \sqrt{8\pi}|Q|f - \frac{2m}{r} + \frac{|Q|^2}{r^2}$) with additional global constant radial electric field (15).

In particular, when $Q = 0$ the solution (17) reduces to electrically neutral black hole with Schwarzschild-anti-de-Sitter geometry for $|\Lambda| > 2\pi f^2$:

$$A(r) = 1 - \frac{2m}{r} + \frac{1}{3}(|\Lambda| - 2\pi f^2)r^2,$$

or electrically neutral black hole with Schwarzschild-de-Sitter geometry for $|\Lambda| < 2\pi f^2$. In both cases above an additional global constant radial electric field (15) is present.

In the special case $|\Lambda| = 2\pi f^2$ and $Q = 0$ we obtain asymptotically flat Schwarzschild black hole ($A(r) = 1 - \frac{2m}{r}$) with additional global constant radial electric field (15), i.e., $\Lambda_{\text{eff}} = 0$ in spite of the presence of the negative bare cosmological constant in the gravity-gauge-field action (16).

4. Charged Test Particle Dynamics

Let us now briefly discuss the dynamics of a test particle with mass $m_0$ and electric charge $q_0$ in the above black hole backgrounds – (6) with (12) and (9) or (6) with (17) and (9). It is given by the standard reparametrization invariant point-particle action:

$$S_{\text{particle}} = \int d\lambda \left[ \frac{1}{2e} g_{\mu\nu}(x) \dot{x}^\mu \dot{x}^\nu - \frac{1}{2} e m_0^2 - q_0 x^\mu A_\mu(x) \right],$$

where $e$ denotes the world-line “einbein”. The standard treatment, using energy $E$ and angular momentum $J$ conservation in the static spherically symmetric background under consideration and replacing the arbitrary world-line parameter $\lambda$ with the particle proper-time parameter $s$ via $\frac{d\lambda}{ds} = em_0$, yields the radial motion equation:

$$\left( \frac{dr}{ds} \right)^2 + V_{\text{eff}}(r) = \frac{E^2}{m_0^2},$$

where $V_{\text{eff}}(r) = \frac{1}{2} (|\Lambda| - 2\pi f^2)r^2$. 


\[ V_{\text{eff}}(r) \equiv A(r) \left( 1 + \frac{J^2}{m_0 r^2} \right) - \frac{q_0^2}{m_0^2} r^2 \left( \frac{\varepsilon_F f}{\sqrt{2}} - \frac{Q}{\sqrt{4\pi r^2}} \right)^2 - 2 \frac{\mathcal{E} q_0}{m_0^2} r \left( \frac{\varepsilon_F f}{\sqrt{2}} - \frac{Q}{\sqrt{4\pi r^2}} \right), \]  

(22)

with \( A(r) \) as in (12) or (17).

Taking for simplicity \( Q = 0 \) (neutral black hole background) and \( J = 0 \) (zero impact parameter – purely radial motion) the “effective” potential (22) becomes:

\[ V_{\text{eff}}^{(0)}(r) = 1 - \frac{2m}{r} + \frac{1}{3} (|\Lambda| - 2\pi f^2) - \frac{q_0^2 f^2}{2m_0^2} r^2 \left( \varepsilon_F f - \frac{Q}{\sqrt{4\pi r^2}} \right). \]  

(23)

In a Schwarzschild-anti-de-Sitter black hole (19) with constant radial electric field (15), for the special value of the ratio of the test particle parameters \( q_0^2/m_0^2 = 2/3(|\Lambda|/f^2 - 2\pi) \) the term quadratic w.r.t. \( r \) in (23) vanishes and the latter acquires the form of a QCD-like (“Cornell”-type [3]) confining-type potential (provided \( q_0 \varepsilon_F < 0 \) with \( \varepsilon_F \) as in (8)):

\[ V_{\text{eff}}^{(0)}(r) = 1 - \frac{2m}{r} + \frac{\sqrt{2} \mathcal{E} q_0 |f|}{m_0^2} r. \]  

(24)

Let us particularly stress that the “Cornell”-type confining potential (24) for charged test particles is exclusively due to the presence of the constant vacuum electric field (15) even though Schwarzschild-anti-de-Sitter is an electrically neutral background.

5. Discussion

It is possible to rewrite the action (2) in an explicitly Weyl-conformally invariant form using the method of two volume forms (two integration measures) [6] introduced earlier in the context of gravity-matter models with primary applications in cosmology. Namely, apart from the standard reparametrization covariant integration density \( \sqrt{-g} \) in terms of the intrinsic Riemannian metric \( g_{\mu\nu} \) as in (2), one introduces an alternative reparametrization covariant integration density \( \Phi(\varphi) \) in terms of auxiliary scalar fields \( \varphi^I \) \((I = 1, \ldots, 4)\):

\[ \Phi(\varphi) = \frac{1}{4!} \varepsilon^{\kappa\lambda\mu\nu} \varepsilon_{IJKL} \partial_{\kappa} \varphi^I \partial_{\lambda} \varphi^J \partial_{\mu} \varphi^K \partial_{\nu} \varphi^L. \]  

(25)
Then the following gravity-gauge-field action:

\[ S = \int d^4x \Phi(\varphi) \left[ g^{\mu\nu} R_{\mu\nu}(\Gamma) \frac{f}{2 \sqrt{-F^2}} - \frac{1}{4} \int d^4x \sqrt{-g} F^2 \right], \]  

where \( R_{\mu\nu}(\Gamma) \) is the Ricci tensor in the first order formalism (i.e., function of the affine connection \( \Gamma^\mu_{\nu\lambda} \)), is explicitly invariant under Weyl-conformal gauge transformations:

\[ g_{\mu\nu} \rightarrow \rho(x) g_{\mu\nu}, \quad \varphi^I \rightarrow \bar{\varphi}^I = \bar{\varphi}^I(\varphi) \text{ such that } \det \left| \frac{\partial \bar{\varphi}^I}{\partial \varphi^J} \right| = \rho(x). \]  

The original action (2) arises as a special gauge-fixed version of the Weyl-conformally invariant action (26) upon using the gauge \( \Phi(\varphi) = \sqrt{-g} \).

To conclude let us recapitulate the main results in the present note:

(a) The non-Maxwell gauge field dynamics of the nonlinear action \( L(F^2) \) in curved space-time (2) produces dynamically a non-zero positive cosmological constant;

(b) The coupled gravity-non-Maxwell-gauge-field system ((2) or (16)) exhibits asymptotically de Sitter and asymptotically non-de Sitter static spherically symmetric (charged) black hole solutions with an additional constant radial electric field (apart from the Coulomb one);

(c) Under certain choice of parameters we find a QCD-like confining-type effective potential in charged test particle dynamics in the above black hole backgrounds.

Furthermore, one can prove the following inverse statement. If we start with an action \( S = \int d^4x \sqrt{-g} \left( R(g) + L(F^2) \right) \) with an a priori unknown gauge field lagrangian \( L(F^2) \) and demand that this theory will possess static spherically symmetric solutions of Reissner-Nordström-de Sitter type, with dynamically generated (via \( L(F^2) \)) cosmological constant, then we derive a unique solution \( L(F^2) = -\frac{1}{4} F^2 - \frac{f^2}{2 \sqrt{-F^2}} \), which was our starting point in (2).

Going back to the non-linear gauge field Eqs.(4) we observe that there exists a more general vacuum solution of the latter without the assumption of staticity and spherical symmetry:

\[ -F^2 = f^2 = \text{const} \quad \text{(recall } F^2 \equiv F_{\kappa\lambda} F_{\mu\nu} g^{\kappa\mu} g^{\lambda\nu}). \]  

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which via Eq.(5) automatically produces an effective positive cosmological constant:

\[ T^{(F)}_{\mu\nu} = -\frac{f^2}{4} g_{\mu\nu}, \quad \text{i.e. } \Lambda_{\text{eff}} = 2\pi f^2. \]  

(29)

Thus, because of the absence of Coulomb field due to (28) and assuming absence of magnetic field, we obtain the above described Schwarzschild-(anti)-de-Sitter (19) or purely Schwarzschild solutions with a vacuum electric field, which according to (28) has constant magnitude \(|\vec{E}| = \sqrt{-\frac{1}{2} F^2} = \frac{f}{\sqrt{2}}\) but its orientation is completely arbitrary. In this disordered vacuum, where the electric field with constant magnitude does not point in one fixed direction, a test charged particle will not be able to get energy from the electric field, instead, it will undergo a kind of brownian motion, therefore no Schwinger pair-creation mechanism will take place.

The present considerations may be extended to higher space-time dimensions and thus provide a framework to study novel effects that could appear relevant in the context of TeV gravity [7] scenarios where nontrivial gauge field effects and gravity effects may be of same order.

As a final comment we mention two other interesting phenomena triggered by the gravity/non-linear-gauge-field system (2) in the context of wormhole physics. First, Misner-Wheeler “charge without charge” effect [8] is known to be one of the most interesting physical phenomena produced by wormholes. Misner and Wheeler realized that wormholes connecting two asymptotically flat space-times provide the possibility of existence of electromagnetically non-trivial solutions, where without being produced by any charge source the flux of the electric field flows from one universe to the other, thus giving the impression of being positively charged in one universe and negatively charged in the other universe.

In an accompanying note [9] we found the opposite effect in wormhole physics, namely, that a genuinely charged matter source of gravity and electromagnetism may appear electrically neutral to an external observer. We show in [9] that this phenomenon takes place when coupling the gravity/gauge-field system (2) self-consistently to a codimension-one charged lightlike brane as a matter source. The “charge-hiding” effect occurs in a self-consistent wormhole solution of the above coupled gravity/gauge-field/lightlike-brane system which connects a non-compact “universe”, comprising the exterior region of Schwarzschild-de-Sitter black hole beyond the internal (Schwarzschild-type horizon), to a Levi-Civita-Bertotti-Robinson-type “universe” with two
compactified dimensions (cf. [10]) via a wormhole “throat” occupied by the charged lightlike brane. In this solution the whole electric flux produced by the charged lightlike brane is expelled into the compactified Levi-Civita-Bertotti-Robinson-type “universe” and, consequently, the brane is detected as neutral by an observer in the Schwarzschild-de-Sitter “universe”.

The above “charge-hiding” solution can be further generalized to a truly charge-confining wormhole solution [11] when we couple the gravity/gauge-field system (2) self-consistently to two separate codimension-one charged lightlike branes with equal but opposite charges. Namely, the latter system possesses a “two-throat” wormhole solution where the “left-most” and the “right-most” “universes” are two identical copies of the exterior region of the neutral Schwarzschild-de-Sitter black hole beyond the Schwarzschild horizon, whereas the “middle” “universe” is of generalized Levi-Civita-Bertotti-Robinson “tube-like” form with geometry $dS_2 \times S^2$ ($dS_2$ being the two-dimensional de Sitter space). It comprises the finite-extent intermediate region of $dS_2$ between its two horizons. Both “throats” are occupied by the two oppositely charged lightlike branes and the whole electric flux produced by the latter is confined entirely within the middle finite-extent “tube-like” “universe”.

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