Author's personal copy

Physics Letters B 704 (2011) 230–233

Contents lists available at SciVerse ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Asymptotically de Sitter and anti-de Sitter black holes with confining electric potential

Eduardo Guendelman [a,*], Alexander Kaganovich [a], Emil Nissimov [b], Svetlana Pacheva [b]

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

ARTICLE INFO

Article history:
Received 2 August 2011
Received in revised form 28 August 2011
Accepted 1 September 2011
Available online 7 September 2011
Editor: M. Cvetič

Keywords:
Black holes in (anti) de Sitter spaces
Dynamically generated cosmological constant
QCD-inspired confining potential

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.

© 2011 Elsevier B.V. All rights reserved.

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]:

$$S = \int d^4x\, L(F^2), \qquad L(F^2) = -\frac{1}{4}F^2 - \frac{f}{2}\sqrt{-F^2},$$
$$F^2 \equiv F_{\mu\nu}F^{\mu\nu}, \qquad F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu, \tag{1}$$

with f being a positive coupling constant. It has been shown in first three references in [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 $\vec{D} = \vec{E} - \frac{f}{\sqrt{2}}\frac{\vec{E}}{|\vec{E}|}$. The pertinent energy density turns out to be (there is *no* contribution from the square-root term in (1)) $\frac{1}{2}\vec{E}^2 = \frac{1}{2}|\vec{D}|^2 + \frac{f}{\sqrt{2}}|\vec{D}| + \frac{1}{4}f^2$, 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 reference in [2].

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

* Corresponding author. Tel.: +972 8 647 2508; fax: +972 8 647 2904.
E-mail addresses: guendel@bgu.ac.il (E. Guendelman), alexk@bgu.ac.il (A. Kaganovich), nissimov@inrne.bas.bg (E. Nissimov), svetlana@inrne.bas.bg (S. Pacheva).
URLs: http://eduardo.hostoi.com (E. Guendelman), http://profiler.bgu.ac.il/frontoffice/ShowUser.aspx?id=1249 (A. Kaganovich), http://theo.inrne.bas.bg/~nissimov/ (E. Nissimov), http://theo.inrne.bas.bg/~svetlana/ (S. Pacheva).

0370-2693/$ – see front matter © 2011 Elsevier B.V. All rights reserved.
doi:10.1016/j.physletb.2011.09.003

Author's personal copy

(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 $G_N = 1$):

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

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:

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = 8\pi T^{(F)}_{\mu\nu}, \tag{3}$$
$$\partial_\nu\left(\sqrt{-g}\left(1 - \frac{f}{\sqrt{-F^2}}\right)F_{\kappa\lambda}g^{\mu\kappa}g^{\nu\lambda}\right) = 0, \tag{4}$$

is explicitly given by:

$$T^{(F)}_{\mu\nu} = \left(1 - \frac{f}{\sqrt{-F^2}}\right)F_{\mu\kappa}F_{\nu\lambda}g^{\kappa\lambda} - \frac{1}{4}\left(F^2 + 2f\sqrt{-F^2}\right)g_{\mu\nu}. \tag{5}$$

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

$$ds^2 = -A(r)\,dt^2 + \frac{dr^2}{A(r)} + r^2\left(d\theta^2 + \sin^2\theta\, d\varphi^2\right), \tag{6}$$
$$F_{\mu\nu} = 0 \quad \text{for } (\mu,\nu) \neq (0,r), \qquad F_{0r} = F_{0r}(r). \tag{7}$$

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

$$\partial_r\left(r^2\left(F_{0r} - \frac{\varepsilon_F f}{\sqrt{2}}\right)\right) = 0, \qquad \varepsilon_F \equiv \mathrm{sign}(F_{0r}), \tag{8}$$

whose solution reads:

$$F_{0r} = \frac{\varepsilon_F f}{\sqrt{2}} + \frac{Q}{\sqrt{4\pi}r^2}, \qquad \varepsilon_F = \mathrm{sign}(Q). \tag{9}$$

Again, as in the flat space–time case (1), the electric field contains 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 $T^0_0 = T^r_r$, which is fulfilled in the present case (7), it is sufficient to solve only one Einstein equation:

$$R^0_0 = 8\pi\left(T^0_0 - \frac{1}{2}T^\lambda_\lambda\right) \quad \text{where } R^0_0 = -\frac{1}{2r^2}\partial_r\left(r^2\partial_r A\right). \tag{10}$$

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

$$8\pi\left({T^{(F)}}^0_0 - \frac{1}{2}{T^{(F)}}^\lambda_\lambda\right) = -4\pi\left(\frac{Q^2}{4\pi r^4} - \frac{1}{2}f^2\right) \tag{11}$$

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:

$$A(r) = 1 - \frac{2m}{r} + \frac{Q^2}{r^2} - \frac{2\pi f^2}{3}r^2. \tag{12}$$

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)):

$$\Lambda_{\mathrm{eff}} = 2\pi f^2. \tag{13}$$

In particular, when $Q = 0$ we obtain electrically neutral black hole with Schwarzschild–de Sitter geometry:

$$A(r) = 1 - \frac{2m}{r} - \frac{2\pi f^2}{3}r^2, \tag{14}$$

where the cosmological constant (13) is *dynamically generated*, and with additional global constant radial electric field:

$$F_{0r} = \varepsilon_F f/\sqrt{2}. \tag{15}$$

3. Bare negative cosmological constant versus induced cosmological constant

Let us now introduce in (2) from the very beginning a negative bare cosmological constant $\Lambda = -|\Lambda|$:

$$S = \int d^4x \sqrt{-g}\left[\frac{1}{16\pi}\left(R(g) - 2\Lambda\right) - \frac{1}{4}F^2 - \frac{f}{2}\sqrt{-F^2}\right]. \tag{16}$$

Then the corresponding static spherically symmetric solution is given by (9) and (6) with:

$$A(r) = 1 - \frac{2m}{r} + \frac{Q^2}{r^2} + \frac{1}{3}\left(|\Lambda| - 2\pi f^2\right)r^2. \tag{17}$$

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:

$$\Lambda_{\mathrm{eff}} = -|\Lambda| + 2\pi f^2 < 0, \quad \textit{i.e.} \quad |\Lambda| > 2\pi f^2. \tag{18}$$

On the other hand, if $|\Lambda| < 2\pi f^2$, *i.e.* $\Lambda_{\mathrm{eff}} = 2\pi f^2 - |\Lambda| > 0$, the solution (17) describes asymptotically de Sitter black hole *in spite* of the presence of negative bare cosmological constant Λ. 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 black hole ($A(r) = 1 - \frac{2m}{r} + \frac{Q^2}{r^2}$) 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 $|\Lambda| > 2\pi f^2$:

$$A(r) = 1 - \frac{2m}{r} + \frac{1}{3}\left(|\Lambda| - 2\pi f^2\right)r^2, \tag{19}$$

Author's personal copy

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.*, $\Lambda_{\rm eff} = 0$ 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:

$$S_{\rm particle} = \int d\lambda \left[\frac{1}{2e} g_{\mu\nu}(x)\dot{x}^\mu \dot{x}^\nu - \frac{1}{2} e m_0^2 - q_0 \dot{x}^\mu A_\mu(x)\right], \qquad (20)$$

where e denotes the world-line "einbein". The standard treatment, using energy $\mathcal{E}$ and angular momentum $\mathcal{J}$ conservation in the static spherically symmetric background under consideration and replacing the arbitrary world-line parameter λ with the particle proper-time parameter s via $\frac{ds}{d\lambda} = e m_0$, yields the radial motion equation:

$$\left(\frac{dr}{ds}\right)^2 + \mathcal{V}_{\rm eff}(r) = \frac{\mathcal{E}^2}{m_0^2}, \qquad (21)$$

$$\mathcal{V}_{\rm eff}(r) \equiv A(r)\left(1 + \frac{\mathcal{J}^2}{m_0 r^2}\right) - \frac{q_0^2}{m_0^2} r^2 \left(\frac{\varepsilon_F f}{\sqrt{2}} - \frac{Q}{\sqrt{4\pi} r^2}\right)^2 - 2\frac{\mathcal{E} q_0}{m_0^2} r \left(\frac{\varepsilon_F f}{\sqrt{2}} - \frac{Q}{\sqrt{4\pi} r^2}\right), \qquad (22)$$

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

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

$$\mathcal{V}^{(0)}_{\rm eff}(r) = 1 - \frac{2m}{r} + \left(\frac{1}{3}\left(|\Lambda| - 2\pi f^2\right) - \frac{q_0^2 f^2}{2m_0^2}\right) r^2 - \frac{\sqrt{2}\mathcal{E} q_0 \varepsilon_F f}{m_0^2} r. \qquad (23)$$

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 ε_F as in (8)):

$$\mathcal{V}^{(0)}_{\rm eff}(r) = 1 - \frac{2m}{r} + \frac{\sqrt{2}\mathcal{E}|q_0| f}{m_0^2} r. \qquad (24)$$

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 φ^I ($I = 1, \ldots, 4$):

$$\Phi(\varphi) = \frac{1}{4!}\varepsilon^{\kappa\lambda\mu\nu}\varepsilon_{IJKL}\partial_\kappa \varphi^I \partial_\lambda \varphi^J \partial_\mu \varphi^K \partial_\nu \varphi^L. \qquad (25)$$

Then the following gravity–gauge-field action:

$$S = \int d^4x\, \Phi(\varphi)\left[\frac{g^{\mu\nu}R_{\mu\nu}(\Gamma)}{16\pi} - \frac{f}{2}\sqrt{-F^2}\right] - \frac{1}{4}\int d^4x \sqrt{-g} F^2, \qquad (26)$$

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:

$$g_{\mu\nu} \to \rho(x) g_{\mu\nu}, \qquad \varphi^I \to \bar{\varphi}^I = \bar{\varphi}^I(\varphi) \quad \text{such that} \quad \det\left\|\frac{\partial \bar{\varphi}^I}{\partial \varphi^J}\right\| = \rho(x). \qquad (27)$$

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 Letter:

(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 *anti*-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 $S = \int d^4x \sqrt{-g}(\frac{R(g)}{16\pi} + L(F^2))$ 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 $L(F^2) = -\frac{1}{4}F^2 - \frac{f}{2}\sqrt{-F^2}$, which was our starting point in (2).

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

$$-F^2 = f^2 = \text{const} \quad \left(\text{recall } F^2 \equiv F_{\kappa\lambda}F_{\mu\nu}g^{\kappa\mu}g^{\lambda\nu}\right), \qquad (28)$$

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

$$T^{(F)}_{\mu\nu} = -\frac{f^2}{4} g_{\mu\nu}, \quad i.e. \quad \Lambda_{\rm eff} = 2\pi f^2. \qquad (29)$$

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

Author's personal copy

field with constant magnitute 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 non-trivial 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".

Acknowledgements

E.N. and S.P. are supported by Bulgarian NSF grant DO 02-257. Also, all of us acknowledge support of our collaboration through the exchange agreement between the Ben-Gurion University of the Negev (Beer-Sheva, Israel) and the Bulgarian Academy of Sciences.

References

[1] G. 't Hooft, Nucl. Phys. B (Proc. Suppl.) 121 (2003) 333, arXiv:hep-th/0208054.
[2] P. Gaete, E. Guendelman, Phys. Lett. B 640 (2006) 201, arXiv:hep-th/0607113;
P. Gaete, E. Guendelman, E. Spalluci, Phys. Lett. B 649 (2007) 217, arXiv:hep-th/0702067;
E. Guendelman, Int. J. Mod. Phys. A 19 (2004) 3255, arXiv:hep-th/0306162;
E. Guendelman, Mod. Phys. Lett. A 22 (2007) 1209, arXiv:hep-th/0703139;
I. Korover, E. Guendelman, Int. J. Mod. Phys. A 24 (2009) 1443;
E. Guendelman, Int. J. Mod. Phys. A 25 (2010) 4195, arXiv:1005.1421 [hep-th].
[3] E. Eichten, K. Gottfried, T. Kinoshita, J. Kogut, K. Lane, T.-M. Yan, Phys. Rev. Lett. 34 (1975) 369;
W. Buchmüller (Ed.), Quarkonia, Current Physics – Sources and Comments, vol. 9, North-Holland, 1992;
M. Karliner, B. Keren-Zur, H. Lipkin, J. Rosner, Ann. of Phys. 324 (2009) 2, arXiv:0804.1575 [hep-ph].
[4] H. Nielsen, P. Olesen, Nucl. Phys. B 57 (1973) 367;
A. Aurilia, A. Smailagic, E. Spallucci, Phys. Rev. D 47 (1993) 2536;
N. Amer, E. Guendelman, Int. J. Mod. Phys. A 15 (2000) 4407.
[5] E. Guendelman, A. Rabinowitz, Gen. Rel. Grav. 28 (1996) 117;
see also T. Jacobson, Class. Quantum Grav. 24 (2007) 5717, arXiv:0707.3272 [gr-qc].
[6] E. Guendelman, A. Kaganovich, Phys. Rev. D 75 (2007) 083505, gr-qc/0607111;
E. Guendelman, Mod. Phys. Lett. A 14 (1999) 1043, gr-qc/0303048;
E. Guendelman, Phys. Lett. B 580 (2004) 87, gr-qc/0303048;
E. Guendelman, A. Kaganovich, Phys. Rev. D 60 (1999) 065004, gr-qc/9905029.
[7] N. Arkani-Hamed, S. Dimopoulos, G. Dvali, Phys. Rev. D 59 (1999) 086004, hep-ph/9807344.
[8] C. Misner, J.A. Wheeler, Ann. of Phys. 2 (1957) 525.
[9] E. Guendelman, A. Kaganovich, E. Nissimov, S. Pacheva, arXiv:1108.3735 [hep-th].
[10] T. Levi-Civita, Rend. R. Acad. Naz. Lincei 26 (1917) 519;
B. Bertotti, Phys. Rev. D 116 (1959) 1331;
I. Robinson, Bull. Akad. Pol. 7 (1959) 351.
[11] E. Guendelman, A. Kaganovich, E. Nissimov, S. Pacheva, arXiv:1109.0453 [hep-th].

Author's personal copy

Physics Letters B 705 (2011) 545

Contents lists available at SciVerse ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Erratum

Erratum to "Asymptotically de Sitter and anti-de Sitter black holes with confining electric potential" [Phys. Lett. B 704 (2011) 230]

Eduardo Guendelman[a,*], Alexander Kaganovich[a], Emil Nissimov[b], Svetlana Pacheva[b]

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

A R T I C L E   I N F O

Article history:
Received 17 October 2011
Accepted 18 October 2011
Available online 20 October 2011
Editor: M. Cvetič

There is a missing second line in Eq. (10) which reads:

$$R^{\theta}_{\theta} = -8\pi T^{0}_{0} \quad \text{where } R^{\theta}_{\theta} = -\frac{1}{r^2}(A-1) - \frac{1}{r}\partial_r A.$$

There is a missing constant term "$-\sqrt{8\pi}|Q|f$" on the right-hand sides of Eqs. (12) and (17).

In the paragraph following Eq. (18) the expression "*asymptotically flat Reissner–Nordström black hole* $(A(r) = 1 - \frac{2m}{r} + \frac{Q^2}{r^2})$" must be replaced by the expression "*flat Reissner–Nordström-like black hole* $(A(r) = 1 - \sqrt{8\pi}|Q|f - \frac{2m}{r} + \frac{Q^2}{r^2})$".

In "Acknowledgement" section an additional sentence should be present: "*We are grateful to Stoycho Yazadjiev for constructive discussions*."

All results remain unchanged.

DOI of original article: 10.1016/j.physletb.2011.09.003.
* Corresponding author.
E-mail address: guendel@bgu.ac.il (E. Guendelman).

0370-2693/$ – see front matter © 2011 Elsevier B.V. All rights reserved.
doi:10.1016/j.physletb.2011.10.042