Quantum Probes of Repulsive Singularities in N = 2 Supergravity

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Abstract

Repulsive singularities ("repulsons") in extended supergravity theories are investigated. These repulsive singularities are related to attractive singularities ("black holes") in moduli space of extended supergravity vacua. In order to study these repulsive singularities a scalar test-particle in the background of a repulson is investigated. It is shown, using a partial wave expansion, that the wave function of the scalar particle vanishes at the curvature singularity at the origin. In addition the connection to higher dimensional p-brane solutions including anti-branes is discussed.

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I. INTRODUCTION

In the last years there has been a lot of progress in understanding black hole physics in supergravity and string theory in $N > 1$ supersymmetric vacua. However, most of the corresponding black hole solutions have a “dual” gravitational repulsive solution in moduli space. Although the existence of these repulsive singularities is known \[1,2\], the repulsive singularities themselves have not been studied much, yet. The main purpose of this article is to consider these gravitational repulsive supersymmetric singularities\[1\] (“repulsons”) and the nature of the corresponding antigravity effects\[2\]. In particular a special class of $N = 2$ supersymmetric BPS repulsons is studied and a general quantum mechanical analysis of a scalar test-particle in the background of these repulsons is given by the use of a partial wave expansion.

The article is organized as follows: In section two static $N = 2$ supergravity solutions preserving $1/2$ of $N = 2$ supersymmetry are introduced. Special points/lines in moduli space where these solutions become either black holes or repulsons are investigated. In section three a scalar test-particle in the background of a repulson is studied. In section four the connection to higher-dimensional $p$-brane solutions is discussed and a conclusion can be found in section five.

II. STATIC SOLUTIONS OF $N = 2$ SUPERGRAVITY

The vector couplings of local $N = 2$ supersymmetric Yang-Mills theory are encoded in the holomorphic function $F(X)$ \[4\], where the $X^I$ ($I = 0 \ldots N_V$) denote the complex scalar fields of the vector supermultiplets. Here $N_V$ counts the number of physical scalars, and $I$ counts the number of physical vectors. The $N_V$ physical scalars parametrize a $N_V$ dimensional complex hypersurface, defined by the condition that the periods satisfy a symplectic constraint. This hypersurface can be described in terms of a complex projective space with

\[1\] In \[1\] these solutions were also called “white” holes, but here we will call them “repulsons” in order to stress that there appears an additional curvature singularity in comparison to the “dual” black holes in moduli space.

\[2\] In general it is known for a long time that antigravity occurs in extended supergravity \[3\].
coordinates $z^A (A = 1, \ldots N_V)$, if the complex coordinates are proportional to some holomorphic sections $X^I(z)$ of the complex projective space: $X^I = e^{K(z, \bar{z})/2} X^I(z)$ with

$$K(z, \bar{z}) = -\log \left( i \bar{X}^I(z) F_I(X^I(z)) - i X^I(z) \bar{F}_I(\bar{X}^I(z)) \right).$$

(II.1)

Moreover one can introduce special coordinates $X^0(z) = 1$ and $X^A(z) = z^A$ with $F(z) = i(X^0)^{-2} F(X)$. It has been shown in [7,11] that the black hole entropy [6] in $N = 2$ supergravity is a topological quantity and given by the central charge $Z$, if the central charge has been minimized [21] with respect to the moduli ($\partial_A |Z| = 0$). Moreover, Behrndt, Lüst and Sabra have shown [5] that one can find general stationary BPS solutions in $N = 2$ supergravity by solving the following $2N_V + 2$ “stabilization equations”

$$X^I - \bar{X}^I = i H^I(r) = i \left( h^I + \frac{p^I}{r} \right)$$

$$F_I - \bar{F}_I = i H_I(r) = i \left( h_I + \frac{q_I}{r} \right)$$

(II.2)

Here the harmonic functions $H(r)$ are given by the constants $h$ and the electric and magnetic charges $q$ and $p$, respectively. The charges satisfy the Dirac quantisation condition

$$p q = 2\pi n, \quad n \in \mathbb{Z}.$$  

(II.3)

Moreover, we restrict ourselves here to static spherically symmetric solutions with metric

$$ds^2 = -e^{-2V(r)} dt^2 + e^{2V(r)} (dr^2 + r^2 d\Omega_2^2)$$

(II.4)

Note that the metric function $e^{2V}$ is given by the holomorphic sections

$$e^{2V} \equiv i \left( \bar{X}^I F_I - X^I \bar{F}_I \right).$$

(II.5)

For black hole solutions the solution of the stabilization equations are on the horizon equivalent to minimizing the central charge with respect to the moduli. Thus, the corresponding black hole entropy [5] is given by

$$S_{BH} = \frac{A}{4G_N} = \lim_{r \to 0} \pi r^2 e^{2V(r)} = \pi |Z_{fix}|^2.$$  

(II.6)

In order to discuss static solutions in $N = 2$ supergravity we have to determine the prepotential. We will consider prepotentials of the form
with intersection numbers $C_{IJK}$. In particular we define special coordinates $(S, T, U) = -i z^{1,2,3}$ and will consider the prepotential

$$F(S, T, U) = -STU - a U^3$$

(II.8)

Here we keep the parameter $a$ to be an arbitrary real number. However, for $a = 1/3$ this prepotential corresponds to the heterotic $S$-$T$-$U$ model with constant hypermultiplets in the large moduli limit. In this model $a$ parametrizes additional perturbative quantum corrections of the prepotential (see [8–10] and reference therein). Moreover, the microscopic interpretation, the higher order curvature corrections and the near-extremal approximation of this class of $N = 2$ models has been studied extensively [19]. For simplicity we take all moduli to be axion-free and $X^0 = \bar{X}^0$. Solving the stabilisation equations with respect to these constraints yields

$$S, T, U = \frac{H^{1,2,3}}{2X^0}, \quad X^0 = \frac{1}{2} \sqrt{-D/H_0}$$

(II.9)

with $D = H^1 H^2 H^3 + a (H^3)^3$. At generic points in moduli space this solution can represent black holes with entropy

$$S_{BH} = 2\pi \sqrt{|q_0| \Delta}.$$  

(II.10)

Here we take $q_0 < 0$ and $\Delta = p_1^1 p_2^2 p_3^3 + a (p^3)^3$. On the other hand it is possible to choose the parameters of the harmonic functions in such a way that the solutions correspond to repulsive singular supersymmetric states [7]. In the following we will study one particular class of these repulsions previously studied by Kallosh and Linde in $N = 4$ supersymmetric string theory [1]. Thus, we consider points/lines in moduli space where the metric function

$$e^{4V(r)} = -4 H_0 D = \sum_{n=0}^{4} \frac{c_n}{r^n}$$

(II.11)

is a polynomial in $1/r$ to second order. Thus, the parameters of the harmonic functions are such that $c_{3,4} = 0$. One particular example corresponds to $a = 0$ and two vanishing charges. Moreover there exist one example with non-vanishing charges, which we will discuss in the following. If
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\[ H^1 H^2 + a(H^3)^2 = c, \quad c = \text{const} \quad (\text{II.12}) \]

i.e. one harmonic function is a function of the other two, the condition \( c_{3,4} = 0 \) is also satisfied. In string theory this might correspond to a gauge symmetry enhancement point/line in moduli space where additional massless states occur. If we take \( h = 1 \) and \( H_0 = -(1 + \frac{q_0}{r}) \) the condition (II.12) reads in terms of the charges

\[ p_1 p_2 + a(p^3)^2 = 0, \quad p^1 + p^2 + 2ap^3 = 0 \quad (\text{II.13}) \]

These conditions can be solved and yield

\[ p^1_{\pm} = -(1 + 2a)p^2 \pm 2p^2\sqrt{a + a^2}. \quad (\text{II.14}) \]

Since all charges are real this yields the bound \( a^2 + a \geq 0 \). If the charges satisfy these equations it follows \( c = 1 + a \). Hence, after a suitable coordinate transformation the metric function becomes

\[ e^{4V(r)} = 1 + \frac{q_0 + p^3}{r} + \frac{q_0 p^3}{r^2} \quad (\text{II.15}) \]

The ADM mass of this solution is \( M = (q_0 + p^3)/4 \geq 0 \). From the Dirac quantisation condition (II.3) follows \( q_0 p^3 = 2\pi n \) with integer \( n \). The solution corresponds to a black hole if \( n \geq 0 \) and to a repulson if \( n < 0 \). Moreover, for \( q_0 = -p^3 \) this solution represents a massless repulson. Note that the entropy of these solutions vanishes, although all harmonic functions and charges are non-vanishing.

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Following [1] we will consider the motion of a scalar test-particle in the background of a repulson, i.e. the metric reads

\[ e^{4V(r)} = 1 + \frac{4M}{r} - \frac{4Z^2}{r^2}, \quad M, Z \geq 0. \quad (\text{III.1}) \]

The corresponding Ricci tensor is given by

\[ R_{tt} = -6 \frac{Mr^3 + 3M^2r - 3r^2Z^2 - 12MrZ^2 + 8Z^2}{(r^2 + 4Mr - 4Z^2)^2} \frac{1}{r^2} \]

\[ R_{rr} = R_{\phi\phi} = R_{\theta\theta} = \frac{6Z^2 - 2Mr}{r^4}. \quad (\text{III.2}) \]
Thus, the metric has one curvature singularity at $r = 0$ and another “naked” curvature singularity at

$$r_0 = 2 \left( \sqrt{M^2 + Z^2} - M \right) \quad \text{(III.3)}$$

A. Classical Analysis

In the classical limit the Newtonian potential $\Phi$ is given by

$$\Phi(r) = -\frac{1}{2} (g_{tt} + 1) = -\frac{M}{r} + \frac{Z^2}{r^2}. \quad \text{(III.4)}$$

The corresponding strength of the gravitational field $\Phi' = \frac{M}{r^2} - \frac{2Z^2}{r^3}$ is gravitational attractive at large distances ($r > r_c$) and gravitational repulsive for $r < r_c$. The critical distance where gravitational repulsion and attraction yield a vanishing net force is given by $r_c = 2Z^2/M$. For massless repulsons the Newtonian potential is always repulsive. Using Hamilton-Jacobi theory (see [12] for example) one can show that a test particle of small mass $m$ ($M >> m$), energy $E$ at $r \to \infty$ and angular momentum $L$ needs the following time to move from $r_1$ to $r_2$

$$t = \int_{r_1}^{r_2} dr \frac{e^{4V}}{\sqrt{E^2 e^{4V} - \frac{L^2}{r^2} - m^2 e^{2V}}} \quad \text{(III.5)}$$

It follows that for $L = 0$ a massive test-particle becomes reflected by the repulson at

$$r_{\min} = \frac{2}{\epsilon} \left( \sqrt{M^2 + \epsilon Z^2} - M \right) > r_0, \quad \epsilon = 1 - \frac{m^4}{E^4} \quad \text{(III.6)}$$

Moreover, for $L \neq 0$ a massless test-particle becomes reflected by the repulson at

$$r_{\min} = 2 \left( \sqrt{M^2 + Z^2 + \frac{L^2}{4E^2} - M} \right) > r_0. \quad \text{(III.7)}$$

Note that for a massless repulson this classical analysis is not valid, since we have chosen the mass of the repulson large, so that the center of mass of the two-body problem is given by the repulson. Moreover, for a massless test-particle in the s-state ($m = L = 0$) a quantum mechanical analysis analogous to [13,14] is necessary.
B. Quantum Mechanical Analysis

We consider a scalar test-particle $\psi$ of mass $m$ satisfying the Klein-Gordon equation in the background of the repulson:

$$\partial_{\mu} \left( \sqrt{-g} g^{\mu\nu} \partial_{\nu} \psi \right) = -m^2 \psi$$  \hspace{1cm} (III.8)

Expanding $\psi$ in partial waves $\psi = e^{-i\omega t} R_{kl}(r) Y_{lm}(\theta, \phi)$ (III.8) becomes the Schrödinger equation

$$\Delta_r \psi + (E - V)\psi = 0$$  \hspace{1cm} (III.9)

with

$$E = m^2 + \omega^2,$$
$$\Delta_r = \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r},$$
$$V(r) = -\frac{4M\omega^2}{r} + \frac{4Z^2\omega^2 + l(l+1)}{r^2}.$$  \hspace{1cm} (III.10)

The Schrödinger equation (III.9) can be solved in the quasi-classical regime, near the singularities $r = 0$ and $r = r_0$ and throughout the entire space-time using standard techniques (see [15] for example). First we will consider (III.9) in the quasi-classical regime for small $r$: The de Broglie wavelength is given by

$$\lambda = \frac{r}{s(s+1)}$$

with $s(s + 1) = 4Z^2\omega^2 + l(l+1)$. The quasi-classical regime requires $\left| \frac{d\lambda}{dr} \right| = \frac{1}{\sqrt{s(s+1)}} \ll 1$. Thus, for large $s$, i.e. large angular momentum of $\psi$ or strong gravitational repulsion with $Z^2 \ll M$, the quasi-classical approximation is valid. In this approximation the radial part of the wave function becomes

$$\psi \sim r^s$$  \hspace{1cm} (III.11)

and has no singular behaviour as one approaches the singularity at $r = r_0$. Thus, the singularity at $r = r_0$ is transparent for massive or massless scalar particles at the quantum level. This result is quite general and does not rely on the partial wave expansion of the scalar test-particle since the Klein-Gordon equation of the scalar test-particle is in general

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3Note that $s$ is quantized in $n$ and $l$ for the solution discussed in section II.

4I thank G. Gibbons for pointing this out to me.
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non-singular at \( r = r_0 \) for the repulson background.

Following [1] we will show now that the scalar test-particles are totally reflected at the singularity at \( r = 0 \). In order to study this effect the ingoing wave function of the scalar must be able to tunnel through the singularity at \( r = r_0 \). Here we “suggest” the following generalization of the metric (II.4) for \( r > 0 \)

\[
    ds^2 = -e^{-4V(r)} |e^{2V(r)}| \, dt^2 + |e^{2V(r)}| \, (dr^2 + r^2 d\Omega^2) .
\]

(III.12)

Thus, at \( r = r_0 \) the determinant of the metric changes its sign. The corresponding “suggested” metric for \( 0 < r < r_0 \) reads

\[
    ds^2_{r < r_0} = |e^{-2V(r)}| \, dt^2 + |e^{2V(r)}| \, (dr^2 + r^2 d\Omega^2)
\]

(III.13)

For small \( r < r_0 \) one can take \( E - V \sim \frac{s(s+1)}{r^2} \). Moreover, from the quasi-classical analysis follows that the radial part of the wave function \( \psi \) can be approximated by a polynomial in \( r \), i.e. \( R_{kl} = A r^n \). Solving the Schrödinger equation (III.9) in this regime yields

\[
    n_{\pm} = \frac{1}{2} \left( \pm \sqrt{1 + 4s(s+1) - 1} \right)
\]

(III.14)

From the quasi-classical analysis follows already that \( n_+ \) is the correct choice for the parameter \( n \). However, we will proof this now in general: One can consider a small region of radius \( \rho \) around the origin and replace \( s(s+1)/r^2 \) by \( s(s+1)/\rho^2 \) in the Schrödinger equation (III.9) for small \( r \). Solving now

\[
    \Delta_r \tilde{R} - \frac{s(s+1)}{\rho^2} \tilde{R} = 0
\]

(III.15)

yields

\[
    \tilde{R}(r) = C \frac{\sin(i\tilde{k}r)}{r}, \quad \tilde{k}^2 = s(s+1)/\rho^2.
\]

(III.16)

Note that \( \tilde{R} \) is finite and vanishing at \( r = 0 \). Now we take as general ansatz \( R(r) = A r^{n_+} + B r^{n_-} \). Since \( R \) and \( \tilde{R} \) and their derivatives must be continuous, it follows

\[
    \partial_r \log(r \tilde{R})|_{r=\rho} = \partial_r \log(r R)|_{r=\rho}
\]

(III.17)

This condition yields

\[
    \frac{B}{A} = -\frac{i\sqrt{s(s+1)} \cot(i\sqrt{s(s+1)}) - n_+ - 1}{i\sqrt{s(s+1)} \cot(i\sqrt{s(s+1)}) - n_- - 1} \rho^{n_+ - n_-}.
\]

(III.18)
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It follows \( \lim_{\rho \to 0} B/A = 0 \). Thus, from

\[
\lim_{\rho \to 0} \frac{1}{A} R(\rho) = \lim_{\rho \to 0} \frac{1}{A} \tilde{R}(\rho) = 0
\]

(III.19)

follows for the radial part of the scalar wave function for small \( r \): \( R(r) = Cr^{n+} \) with some constant \( C \). Hence, the scalar wave function vanishes at the singularity \( r = 0 \). This result is valid for massive and massless repulsons.

To investigate the space-time geometry near the reflecting singularity we consider the metric (III.13) and perform a Weyl-transformation of the metric \( g_{\mu\nu} \to |e^{2V}| g_{\mu\nu} \) and an additional coordinate transformation \( r' = -2Z \log r \). This yields the following metric near the singularity \( r = 0 \).

\[
ds^2 = dt^2 + dr'^2 + 4Z^2 d\Omega_2^2.
\]

(III.20)

Thus, the four-dimensional space-time \( M_4 \) near the repelling singularity is in this particular “frame” a product space \( M_4 = \mathbb{R}^2 \times S^2 \), where the radius of \( S^2 \) is associated to the strength of the gravitational repulsion. It follows for the solution discussed in section II that the radius of \( S^2 \) and the corresponding “area” \( A = 16\pi Z^2 \) are quantized, since \( 4Z^2 = 2\pi|n| \).

Now we will solve the Schrödinger equation (III.9) in the region \( r > r_0 \). Introducing the wave number \( k = \sqrt{E} \) and \( \alpha = 2M\omega^2 \) the radial part of the wave function is given by [15]

\[
R_{kl}(r) = \frac{C_k}{(2s+1)!} (2\pi)^{1/2} e^{-ikr} F\left(\frac{i}{k} + s + 1, 2s + 2, \alpha ikr\right)
\]

(III.21)

Here \( F(a, b, x) \) denotes the confluent hypergeometric function and the constant \( C_k \) is defined as follows

\[
C_k = \sqrt{\frac{2}{\pi}} k e^{\frac{\alpha}{2k}} \left| \Gamma\left(s + 1 - i\frac{\alpha}{2k}\right) \right|, \quad \int dr \ r^2 R_{kl} R_{kl'} = 2\pi \delta(k - k').
\]

(III.22)

For \( r \to \infty \) this solution becomes

\[
R_{kl}(r) = \sqrt{\frac{2}{\pi}} \frac{1}{r} \sin(kr + \frac{\alpha}{2k} \log(2kr) - \frac{s\pi}{2} + \delta_s), \quad \delta_s = \arg \left( \Gamma(s + 1 - i\frac{\alpha}{2k}) \right).
\]

(III.23)

This result implies that the scattering process of a scalar particle of mass \( m \) and a heavy repulson of mass \( M >> m \) is elastic, only, and the inelastic cross-section vanishes. Hence, in a scattering process the repulson reflects any scalar test-particle. Note that this conclusion holds for a massive and massless scalar test-particle, but only for a massive and heavy repulson.

\[\text{\footnotesize \( F(a, b, x) = \sum_{n=0}^{\infty} a_n \frac{x^n}{b^n/n!} \) with } a_0 = 1, a_n = a(a+1) \ldots (a+n-1), b_n = b(b+1) \ldots (b+n-1). \]
IV. REPUlSONS FROM EXTREMAL P-BRANES

The class of repulson solutions we have studied so far can be associated to higher dimensional brane configurations. In general they correspond to orthogonally intersecting brane-anti-brane configurations. These configurations become massless if the absolute values of the corresponding two brane-charges are equal. If we consider, for example, four-dimensional configurations in \( N \geq 2 \) supergravity with a prepotential of the form (II.8) with \( a = 0 \) and non-vanishing charges, then the particular solution can be interpreted as the intersection of three 4-branes and one 0-brane in \( D = 10 \) type IIA string theory. In M-theory this solution correspond to the intersection of three 5-branes with a boost along the common string \([10,5]\). Wrapping the 5-branes around 4-cycles this configuration corresponds to a magnetic string in five dimensions (see e.g. \([17,9]\) and reference therein). Additional wrapping around the 5th direction yields the four dimensional solution. The appearance of negative charges defines the corresponding anti-brane configurations and yields massless solutions at particular points in moduli/charge space \([2,5,9]\). Note that solutions with anti-brane intersections, only, have negative ADM mass (see also \([22]\)). Switching off two charges corresponds to two vanishing branes and studying the corresponding brane-anti-brane intersection in four dimensions yields the repulson background studied above \([18]\).

V. SUMMARY, DISCUSSION AND CONCLUSION

Supersymmetric BPS solutions of \( N = 2 \) supergravity have been studied. In particular black holes and repulsons have been considered and it has been shown that both types of singularities are closely related to each other, i.e. a black holes solution at one point in moduli space becomes a repulson at another point \([2]\). For a two-charge configuration, for example, with a generic choice for the moduli at infinity and positive greater charge, the lower charge with a plus sign corresponds to a black hole solution and the lower charge with a minus sign to a repulson. It might be interesting to investigate whether there are as many repulsive as attractive solutions in moduli space, i.e. almost all attractive black hole solutions might have “dual” repulsive solutions. Since a repulson is gravitational attractive at long distances and gravitational repulsive at short distances, it is particle-like \([20,13,18]\). In order to investigate the reflecting nature of a repulson a scalar test-particle in the background of a repulson has been studied \([1]\). It has been shown in general that the wave function of the
scalar vanishes at the singularity $r = 0$ and is transparent at the singularity at $r = r_0$, if the wave function can be expanded in partial waves. Moreover, any scattering process of a heavy repulson and a scalar is elastic, only. These results suffer from the boundary condition at the naked singularity at $r = r_0$. Here we “suggested” a generalization of the canonical metric valid for $r > 0$ to study tunneling of the ingoing scalar test-particle through the naked singularity at $r = r_0$. It would be very interesting to investigate this point further. Moreover, massless repulsions [2,9] are rather special and also deserve further investigations. Note that the second part of this article can be applied to the particular string model with non-Abelian gauge fields studied by Kallosh and Linde in [1] and generalizes their quantum mechanical analysis of a scalar test-particle in the background of a repulson.

From the brane point of view the class of repulsions studied here correspond to orthogonally brane-anti-brane intersections. The corresponding “dual” black hole solution in moduli space would correspond to the brane-brane intersection. Moreover, the corresponding anti-brane-anti-brane intersection is also gravitational attractive, but its ADM mass is negative. In general the four dimensional solutions are gravitational repulsive if they contain an odd number of anti-branes and gravitational attractive if the number of anti-branes is even. To conclude, naked repulsive singularities appearing in extended supergravity models have been studied in great detail. These repulsive singularities are related to black hole solutions in moduli space [2], i.e. in moduli space these repulsive singular solutions are as generic as their attractive “duals”. It would be very interesting to investigate more general repulsive singular solutions with, for instance, additional $1/r^3$ and $1/r^4$ contributions to the metric function or without using a partial wave expansion.

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