Gravitinos tunneling from black holes

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A B S T R A C T

Black hole radiation of gravitinos is investigated as the classically forbidden tunneling of spin-3/2 fermions through an event horizon. We calculate directly that all four spin states of the gravitino yield the same emission temperature, and the Unruh temperature in a Rindler spacetime as well as the Hawking temperature for a Kerr–Newman charged rotating black hole are retrieved. This confirms the robustness of the tunneling formalism in a wide range of applications.

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1. Introduction

In recent years, theorists have begun modeling Hawking radiation as a tunneling effect, offering an alternate conceptual understanding to black hole radiation. Originating with the early work of Kraus and Wilczek [1], the tunneling approach was later refined [2–5] to show that black hole radiation can indeed be understood in terms of local physics near the horizon. This approach is quite distinct from more global methods such as Wick rotation methods and Hawking’s original method of modeling gravitational collapse [6]. Although the approach was first demonstrated for scalar particles tunneling out of spherically symmetric Schwarzschild black holes, it has been shown to be remarkably versatile, and has been applied successfully to a wide variety of spacetimes, including Kerr and Kerr–Newmann cases [7–9], black rings [10], the 3-dimensional BTZ black hole [11,12], Vaidya [13], other dynamical black holes [14], Taub–NUT spacetimes [9], Gödel spacetimes [15], and non-commutative geometries [16]. One can even recover the Unruh temperature [17] for a Rindler observer [4,5,9] using this approach.

The tunneling method has been implemented in two different ways: the null geodesic method, in which a shell of null radiation is emitted from a black hole via a tunneling process [3] and the Hamilton–Jacobi method [11], which is an extension of the complex path approach [4]. All tunneling approaches make use of the WKB approximation, whose validity is predicated on the assertion that gravitational blueshift near the horizon ensures the radiation has a wavelength much shorter than the width of the barrier induced by the horizon, through which the radiation tunnels. Unlike field-theoretic methods that rely on a decomposition of a field into positive and negative energy components (that in turn are related by Bogoliubov transformations, whose coefficients are obtained by requiring that the modes of fixed Killing frequency are analytic when prolonged across the horizon), no explicit identification of positivity of energy is made in tunneling methods. Rather this identification is implicit in the spacetime coordinatization in conjunction with the ansatz used for the calculation of the imaginary part of the action. This identification can be made coordinate-independent through introduction of orthonormal frames [18].

In the present Letter we will consider only the Hamilton–Jacobi method, in which the wavefunction of scalar radiation is approximated to leading order in $\hbar$ by the exponential of an action functional $I$ that satisfies the Hamilton–Jacobi equation. From this the tunneling probability $\Gamma$, given by the ratio of emission to absorption probabilities, is to leading order in $\hbar$ (which we set to unity) [1,19]

$$\Gamma \propto \exp(-2\Im I).$$

Comparing this to a thermal Boltzmann distribution for emission, $\Gamma = \exp(-\beta E)$, allows one to assign a temperature $\beta^{-1}$ to the emitted radiation.

One expects that a black hole should radiate all types of particles, akin to a black body at a well-defined temperature (ignoring grey body effects), such that particles of every spin should appear in the emission spectrum. While the implications of this were studied 30 years ago [20], it has only recently been shown that spin-1/2 fermions can tunnel out of black holes [21]. Although implementation of the Hamilton–Jacobi approach (to leading order in a WKB approximation) is quite distinct from the scalar case, one obtains the same temperature as for scalar radiation. As with the scalar case, the fermionic tunnelling approach has also been shown to be quite robust, and has been applied to a broad variety of black hole spacetimes [21–23].

In this Letter we consider radiation of spin-3/2 fermions, or gravitinos, which are predicted to exist in all theories of supergravity [24]. Since local supersymmetry is not a symmetry of the
state of our observable universe, it should be broken, and so the gravitino should be massive. Extending the tunneling approach to the spin-3/2 case yields additional complications, since the gravitino obeys the Rarita–Schwinger equation [25] and so any one of four possible spin states can be emitted by the black hole. With reference to the radial direction, these components are 3/2, 1/2, −1/2, and −3/2. Since it is not a priori clear that all components are emitted in the same way, we have a new and interesting check on the tunneling method, which we will use to directly calculate the emission probabilities, without resorting to thermodynamic arguments.

Our approach will be analogous to that taken in the spin-1/2 case [21,26]. Employing a WKB approximation to the Rarita–Schwinger equation, we will show that a general structure emerges that ensures all four spin states are radiated by the black hole at the same Hawking temperature. To keep the discussion more general, we will couple the gravitino to a U(1) gauge field as well as gravity. We apply these results first to Rindler spacetime, confirming that the Unruh temperature is recovered. We then consider the Kerr–Newman case, which had introduced some non-trivial technical features associated with the choice of γ matrices for the spin-1/2 case [26]. We show that the same approach is sufficient in describing gravitino emission to leading order in the WKB approximation.

One of the assumptions of our semi-classical calculation is to neglect any change of angular momentum of the black hole due to the spin of the emitted particle. For zero angular momentum black holes with mass much larger than the Planck mass this is a good approximation. Furthermore, particles of opposite spin will be emitted in equal numbers statistically, yielding no net change in angular momentum of the black hole. Thus, since we are only considering a good approximation, we will neglect any change of angular momentum of the black hole due to the spin of the emitted particle. For zero angular momentum black holes with mass much larger than the Planck mass this is a good approximation. Furthermore, particles of opposite spin will be emitted in equal numbers statistically, yielding no net change in the angular momentum of the black hole (although second-order statistical fluctuations will be present). We confirm that spin-3/2 fermions are also emitted at the Hawking temperature. This final result, while not surprising, furnishes an important confirmation of the robustness of the tunneling approach.

2. Gravitinos in a black hole background

The Rarita–Schwinger equation representing the spin-3/2 fermion field will be used in the form [25]

\[ i\hbar \gamma^\nu (D_\nu + i q A_\nu) \Psi_\mu + m \Psi_\mu = 0, \]  \hspace{1cm} (2a)  
\[ \gamma^{\mu
u} \Psi_\mu = 0, \]  \hspace{1cm} (2b)  
where \( A_\nu \) is a vector-valued spinor of charge \( q \) and mass \( m \). \( \gamma_\mu \) represents the electromagnetic potential of the black hole, and the \( \gamma^\mu \) matrices satisfy \( \{ \gamma^\mu, \gamma^\nu \} = 2g^{\mu\nu} \). The covariant derivative obeys

\[ D_\mu = \partial_\mu + \Omega_\mu, \]
\[ \Omega_\mu = \frac{1}{2} \Gamma^{\alpha\beta}_{\mu} \gamma_{\alpha\beta}, \]
\[ \Gamma^\mu_{\alpha\beta} = \frac{1}{4} [\gamma^\mu, \gamma^{\alpha\beta}], \]  \hspace{1cm} (3)
where \( \Omega_\mu \) is the spin-connection.

The first equation is the Dirac equation applied to every vector index of \( \Psi \), while the second is a set of additional constraints to ensure that no ghost state propagates; that is, to ensure that \( \Psi \) represents only spin-3/2 fermions, with no spin-1/2 mixed states.

Working in the context of the path-integral formalism and the WKB approximation, it is known that each path will have a phase of \( \exp(i l / \hbar) \), where \( l \) is the action corresponding to that path and may depend on the spin eigenstate of the particle. In the case that we consider here, the infinitesimal radial path across the horizon will dominate, such that we may employ the following ansatz for the wave function:

\[ \Psi_\mu = \left[ \begin{array}{c} a_\mu \\ b_\mu \\ c_\mu \\ d_\mu \end{array} \right] \exp i l, \]  \hspace{1cm} (4)
where the \( a_\mu, b_\mu, c_\mu, d_\mu \) are each functions of the spacetime.

Note that if we do not take Rarita–Schwinger equations (2a) we will yield an equation which can be solved for the action \( l \) independently of the wave function components \( a_\mu, \ldots, d_\mu \). The second set of Rarita–Schwinger equations (2b), on the other hand, will yield four constraints for these wave function components independently of the action, and will have solutions in every spacetime. Hence, as the action is all that we require to find the emission temperature, we need solve only Eq. (2a).

This very simple observation has the dramatic and immediate consequence that every spin-3/2 fermion, in any spin state, will be emitted from a black hole at the same temperature. Indeed, choosing a spin state for our wavefunction would be equivalent to choosing functions \( a_\mu, \ldots, d_\mu \) that correspond to the spin eigenstate, allowing us to calculate the action corresponding to this spin state. However, as we have found, those equations will be independent of the functions \( a_\mu, \ldots, d_\mu \), and hence independent of the spin eigenstate. This comes out very naturally and simply, and is the generalization of the known fact that spin-1/2 fermions are emitted at the same temperature regardless of spin state, which is explicitly calculated in [21].

Defining \( \gamma^\mu \) to be of the chiral form,

\[ \gamma^0 = \left[ \begin{array}{cc} 0 & 1 \\ -1 & 0 \end{array} \right], \quad \gamma^i = \left[ \begin{array}{cc} 0 & \sigma^i \\ \sigma^i & 0 \end{array} \right], \]  \hspace{1cm} (5)
where then, in (2a), \( \gamma^\mu \equiv \gamma^\mu_\nu \gamma^\nu \) are metric-dependent linear combinations of these matrices, we find that the Rarita–Schwinger equation (2a) can be approximated to first-order in \( \hbar \) (which is from now on set to unity) and rewritten in the form

\[ \left[ \begin{array}{cccc} m & x_0 + x_3 & x_1 - ix_2 & 0 \\ -x_0 + x_3 & x_1 + ix_2 & x_0 - x_3 & m \\ x_1 + ix_2 & -x_0 - x_3 & m & m \end{array} \right] \left[ \begin{array}{c} a \\ b \\ c \\ d \end{array} \right] = 0, \]  \hspace{1cm} (6)
where

\[ x_a = \epsilon^a_\mu \partial_\mu I. \]  \hspace{1cm} (7)

Performing an LU-decomposition of the matrix above yields a solution independent of the wave function components:

\[ \eta^{ab} (\epsilon^a_\mu \partial_\mu I) (\epsilon^b_\nu \partial_\nu I) - m^2 = 0, \]  \hspace{1cm} (8)
where \( \eta = \text{diag}(-1, 1, 1) \) is the Minkowski flat-space metric. This equation is equivalent to the first set of Rarita–Schwinger equations (2a), and is the same as the coordinate-free formulation studied for the scalar radiation [18]. From it we can solve for \( \partial_\mu I \), regardless of what \( \Psi \) looks like, and hence independently of (2b). Integrating \( \partial_\mu I \) along the infinitesimal radial path across the event horizon (or more properly in the upper-right complex-\( r \) plane) will yield a complex residue corresponding to the action of the radiation. From this action, the temperature is found using [1,19]

\[ \Gamma \propto \exp(-2i m) \propto \exp \left( \frac{-E}{\hbar \sqrt{I}} \right), \]  \hspace{1cm} (9)
Note that the above relations hold regardless of the spin-state chosen for gravitino emission. Consequently, all four spin states are radiated with the same Hawking temperature to leading order in WKB. We will now explicitly demonstrate this method in Rindler and Kerr–Newman spacetimes.

3. Rindler spacetime

We first make use of the results of the preceding section in Rindler spacetime, whose simplicity allows one to easily grasp the concept of the method employed. Although this is not a black hole spacetime, we will show that the tunneling approach allows us to calculate the tunneling probabilities of spin-3/2 particles through the Rindler horizon, yielding the well-known Unruh temperature [17] as for lower-spin particles [3,9,11,21]. This result illustrates the versatility and power of the tunneling approach, linking the Unruh effect to black hole radiation. Consider the metric

\[ ds^2 = -f(z)dt^2 + dx^2 + dy^2 + \frac{dz^2}{g(z)}, \]

where

\[ f(z) = \alpha^2 z^2 - 1, \]

\[ g(z) = \alpha^2 z^2 - 1, \]

\[ A_\mu = 0, \]

which has a singularity at \( z = \alpha^{-1}, \) and consider the action to have a solution to (8) of the form

\[ I = -Et + W(z) + P(x, y) + K. \]

Inserting (13) into the Rarita–Schwinger equation (2a), we obtain, to leading order in \( h \) such that neither \( \Omega_\mu \) nor the derivatives of the \( a_\mu \), \( d_\mu \), contribute, the four equations:

\[ c_\mu \left( \frac{E}{\sqrt{f}} - \sqrt{g} \right) - c_\mu (P_x - iP_y) + a_\mu m = 0, \]

\[ d_\mu \left( \frac{E}{\sqrt{f}} + \sqrt{g} \right) - c_\mu (P_x + iP_y) + b_\mu m = 0, \]

\[ a_\mu \left( \frac{E}{\sqrt{g}} + \sqrt{f} \right) + b_\mu (P_x - iP_y) - c_\mu m = 0, \]

\[ b_\mu \left( \frac{E}{\sqrt{g}} - \sqrt{f} \right) + a_\mu (P_x + iP_y) - b_\mu m = 0. \]

The constraints (2b) will also give us the relations

\[ \frac{c_1}{\sqrt{f}} + d_2 + id_3 + \sqrt{g} c_4 = 0, \]

\[ \frac{d_1}{\sqrt{f}} + c_2 + ic_3 - \sqrt{g} d_4 = 0, \]

\[-a_1 \frac{b_1}{\sqrt{f}} + b_2 - ib_3 + \sqrt{g} a_4 = 0, \]

\[-b_1 \frac{a_1}{\sqrt{f}} + a_2 + ia_3 - \sqrt{g} b_4 = 0. \]

Solving Eqs. (14) exactly and generally, we get the same solution as (8):

\[ \left( \frac{E}{\sqrt{f}} - W' \sqrt{g} \right) \left( \frac{E}{\sqrt{f}} + W' \sqrt{g} \right) - (P_x - iP_y)(P_x + iP_y) + m^2 = 0. \]

As we approach the horizon, we find that \( f(z) \) and \( g(z) \) both go to zero, while the other terms (besides \( W' \)) are constant, and therefore get

\[ W'_\pm = \pm \frac{E}{\sqrt{f}}. \]

This solution will be valid for any non-zero \( \Psi \) satisfying (15), whose solution space generally is 12-dimensional.

We then integrate (17) over our path through the event horizon, which contains a pole that will be the sole imaginary contribution to the action. This yields:

\[ \text{Im } W_{\pm}(z) = \frac{\pm \pi E}{\sqrt{f(z_0)g(z_0)}} = \frac{\pi E}{2a}. \]

As \( W' \) corresponds to the momentum of the particle, one finds that \( W_+ \) corresponds to an outgoing particle, whereas \( W_- \) corresponds to an incoming one. By forcing to unity the probability that the incoming particle is absorbed and using the fact that \( W_+ = -W_- \), we find

\[ \Gamma \propto \exp(-2|E|) = \frac{\exp(-2(\text{Im } W_+ + \text{Im } P(x, y) + \text{Im } K))}{\exp(-2(\text{Im } W_- + \text{Im } P(x, y) + \text{Im } K))} = \exp(-4|\text{Im } W_+|) = \exp \left( -\frac{2\pi |E|}{a} \right), \]

which gives us the expected Unruh temperature [17]:

\[ T_\mu = \frac{\alpha}{2\pi}. \]

4. Kerr–Newman black hole spacetime

We next consider applying this method to a general Kerr–Newman black hole spacetime. The metric is given by

\[ ds^2 = -f dt^2 + \frac{dr^2}{g} - 2H dt d\phi + K d\phi^2 + \Sigma d\theta^2, \]

where

\[ A_\mu = -\frac{\epsilon r}{\Sigma(r)} ((dt)_\mu - \alpha \sin^2(\theta)(d\phi)_\mu), \]

\[ f(r, \theta) = \frac{\Delta(r) - \alpha^2 \sin^2(\theta)}{\Sigma(r, \theta)}, \]

\[ g(r, \theta) = \frac{\Delta(r)}{\Sigma(r, \theta)}, \]

\[ H(r, \theta) = \frac{\alpha \sin^2(\theta)(r^2 + \alpha^2 - \Delta(r))}{\Sigma(r, \theta)}, \]

\[ K(r, \theta) = \frac{(r^2 + \alpha^2)^2 - \Delta(r)\alpha^2 \sin^2(\theta)}{\Sigma(r, \theta)} \sin^2(\theta), \]

\[ \Sigma(r, \theta) = r^2 + \alpha^2 \cos^2(\theta), \]

\[ \Delta(r) = r^2 + \alpha^2 + e^2 - 2Mr = (r - r_-)(r - r_+). \]

We assume a non-extremal black hole, \( M^2 > \alpha^2 + e^2 \), such that we have two horizons at
\[ r_\pm = M \pm \sqrt{M^2 - \alpha^2 - \varepsilon^2}. \]

To simplify the notation, we will also use the functions
\[
F(r, \theta) = f(r, \theta) + \frac{H^2(r, \theta)}{K(r, \theta)} = \frac{\Delta(r)\Sigma\left(\frac{\pi}{2}\right)}{(i + \alpha^2)^2 - \Delta(r)\alpha^2 \sin^2(\theta)}, \tag{24a}
\]
\[\Omega_H = \frac{H(r_+, \theta)}{K(r_+, \theta)} = \frac{a}{r_+^2 + \alpha^2}, \tag{24b}\]

where \(\Omega_H\) corresponds to the angular velocity of the black hole.

We will use the following \(\gamma^{\mu\nu}\) matrices [26]:
\[
\gamma^i = \frac{1}{\sqrt{F(r, \theta)}} \gamma^0, \quad \gamma^\prime = \sqrt{g(r, \theta)} \gamma^3, \quad \gamma^\theta = \frac{1}{\sqrt{\Sigma(r, \theta)}} \gamma^4.
\]

\[
\gamma^\phi = \left(\hat{\Sigma}(r, \theta) + \frac{H(r, \theta)}{\sqrt{F(r, \theta)K(r, \theta)}} \gamma^0\right), \tag{25}\]

where the \(\hat{\gamma}^i\) are the chiral Minkowski matrices (5). The action will in this case take the form
\[
I = -Et + W(r, \theta) + J\phi. \tag{26}\]

If we then input our wave function (4) into the Rarita–Schwinger equations (2a), we get, again to leading order in \(\hbar\), such that neither \(\Omega_\mu\), nor the derivitives of the \(a_\mu, \ldots, d_\mu\) contribute, equations similar to those found for the Rindler spacetime:
\[
c(\zeta - W_r \sqrt{\Sigma}) + d\left(\frac{\xi - W_\theta}{\sqrt{\Sigma}}\right) + am = 0, \tag{27a}\]
\[
d(\zeta + W_r \sqrt{\Sigma}) + c\left(\frac{-\xi - W_\theta}{\sqrt{\Sigma}}\right) + bm = 0, \tag{27b}\]
\[
a(\zeta + W_r \sqrt{\Sigma}) - b\left(\frac{-\xi - W_\theta}{\sqrt{\Sigma}}\right) - cm = 0, \tag{27c}\]
\[
b(\zeta - W_r \sqrt{\Sigma}) - a\left(\frac{-\xi - W_\theta}{\sqrt{\Sigma}}\right) - dm = 0, \tag{27d}\]

where
\[
\zeta = E - \frac{\sqrt{\alpha^2}}{\sqrt{r}} \left(J + \frac{\sqrt{\alpha^2}}{\sqrt{K}} a \sin^2(\theta)\right), \tag{28a}\]
\[
\xi = i \left(\frac{J}{\sqrt{K}} + \frac{\sqrt{\alpha^2}}{\sqrt{K}} a \sin^2(\theta)\right). \tag{28b}\]

Once again, the second set of Rarita–Schwinger equations (2b) will give us additional relations between the various vector components of the wave function, which are not important here as our solution for the action will be independent of such relations:
\[
\frac{c_1}{\sqrt{F}} + \sqrt{\Sigma} c_2 + \frac{d_1}{\sqrt{F}} + \frac{Hc_4}{K\sqrt{F}} - \frac{id_4}{\sqrt{K}} = 0, \tag{29a}\]
\[
\frac{d_1}{\sqrt{F}} - \sqrt{\Sigma} d_2 + \frac{c_1}{\sqrt{F}} + \frac{Hd_4}{K\sqrt{F}} + \frac{ic_4}{\sqrt{K}} = 0, \tag{29b}\]
\[
-a_1 + \sqrt{\Sigma} a_2 + \frac{b_1}{\sqrt{K}} = \frac{Ha_4}{K\sqrt{F}} - \frac{ib_4}{\sqrt{K}} = 0, \tag{29c}\]
\[
-b_1 - \sqrt{\Sigma} b_2 + \frac{a_1}{\sqrt{K}} = \frac{Hb_4}{K\sqrt{F}} + \frac{ia_4}{\sqrt{K}} = 0. \tag{29d}\]

Eqs. (27) may be rewritten as
\[
(\zeta - W_r \sqrt{\Sigma})(\zeta + W_r \sqrt{\Sigma}) - \left(\frac{W_\theta}{\sqrt{\Sigma}} - \xi\right)\left(\frac{W_\theta}{\sqrt{\Sigma}} + \xi\right) + m^2 = 0. \tag{30}\]

Since the action may be reduced to an infinitesimal path from inside to outside the horizon, \(\theta\) may only take values between \(\theta_0 - \varepsilon\) and \(\theta_0 + \varepsilon\). Hence, we make the approximation \(\theta = \theta_0\), which in turns forces \(W_\theta\) to be constant. Then, expanding (30) near the horizon \(r \to r_+\) and solving for \(W_r\), we see that only \(\xi\) contributes, as it rapidly increases to infinity near the horizon while other terms do not. This allows us to find a solution similar to that for emission of Dirac particles [21]:
\[
W_{r \pm} = \pm \sqrt{\frac{2}{r_0}} \frac{\Xi}{\sqrt{\frac{r_0}{r_+}}} = \pm \frac{E - J\Omega_H + \frac{qer_+}{r_+^2 + \alpha^2}}{\sqrt{F(r_+\theta)g(r_+\theta)(r - r_+)}}, \tag{31}\]

Integrating (31) yields
\[
\Gamma = \frac{\pm \pi i (E - J\Omega_H + \frac{qer_+}{r_+^2 + \alpha^2})}{2r_+ - 2M}, \tag{32}\]

which is independent of the particular choice of \(\theta_0\). Thus, taking again \(\Gamma \propto \exp(-4\Im W_{r \pm})\), we find
\[
\Gamma \propto \exp(-2\pi i \frac{r_+^2 + \alpha^2}{2r_+ - 2M} (E - J\Omega_H + \frac{qer_+}{r_+^2 + \alpha^2})), \tag{33}\]

giving us the expected Hawking temperature for a charged rotating black hole,
\[
T_H = \frac{1}{2\pi} \frac{r_+ - M}{r_+^2 + \alpha^2} = \frac{1}{2\pi} \frac{\sqrt{M_+^2 - \alpha^2 - \varepsilon^2}}{2M(M + \sqrt{M^2 - \alpha^2 - \varepsilon^2}) - \varepsilon^2}. \tag{34}\]

5. Conclusions

We have shown that gravitinos, as described by the Rarita–Schwinger equation in a curved spacetime, can tunnel out of event horizons, and have carried out explicit calculations with regards to their emission temperatures. The tunnelling formalism applies equally well to this situation as to previous work on scalar particles and spin-1/2 fermions because the additional constraints from the Rarita–Schwinger equation do not modify the resultant Hamilton–Jacobi equation to leading order in \(\hbar\). The implication is that fermions of every spin will share a common Hawking temperature. We found that gravitinos attain the familiar Unruh temperature in a Rindler spacetime, and recovered the Hawking temperature for gravitino radiation from a Kerr–Newman black hole. All four spin components yield the same temperature to leading order in a WKB approximation. These results are in agreement with those obtained by applying this method to spin-1/2 fermions [21,26], as well as with results obtained by using the Newman–Penrose approach on gravitinos [27] radiating from a Reissner–Nordström black hole.

An interesting task would be to evaluate the corrections to the emission probability that are subleading in \(\hbar\). These could in principle involve mixing between different gravitino components, whose effect on the resultant thermodynamic behaviour could lead to deviations from black body behaviour.

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