Hindawi
Advances in High Energy Physics
Volume 2018, Article ID 8504894, 7 pages
https://doi.org/10.1155/2018/8504894

Research Article

Absorption Cross Section and Decay Rate of Dilatonic Black Strings

Huriye Gürsel and İzzet Sakallı

Physics Department, Faculty of Arts and Sciences, Eastern Mediterranean University, Famagusta, North Cyprus, Mersin 10, Turkey

Correspondence should be addressed to İzzet Sakallı; izzet.sakalli@emu.edu.tr

Received 25 October 2018; Accepted 4 December 2018; Published 19 December 2018

Academic Editor: Piero Nicolini

We studied in detail the propagation of a massive tachyonic scalar field in the background of a five-dimensional (5D) Einstein–Yang–Mills–Born–Infeld–dilaton black string: the massive Klein–Gordon equation was solved, exactly. Next we obtained complete analytical expressions for the greybody factor, absorption cross section, and decay rate for the tachyonic scalar field in the geometry under consideration. The behaviors of the obtained results are graphically represented for different values of the theory's free parameters. We also discuss why tachyons should be used instead of ordinary particles for the analytical derivation of the greybody factor of the dilatonic 5D black string.

1. Introduction

A wealth of information about quantum gravity can be obtained by studying the unique and fascinating objects known as black holes (BHs). In BH physics, greybody factors (GFs) modify black-body radiation, or predicted Hawking radiation [1, 2], within the limits of geometrical optics [3]. In other words, GFs modify the Hawking radiation spectrum observed at spatial infinity (SI), so that the radiation is not purely Planckian [4].

GF, absorption cross section (ACS), and decay rate (DR) are quantities dependent upon both the frequency of radiation and the geometry of spacetime. Currently, although there are many studies of GF, ACS, and DR (see, for example, [5–10] and the references therein), the number of analytical studies of GFs that consider modified black-body radiation of higher-dimensional ($D > 4$) spacetimes, like the BHs in string theory and black strings [11–13], is rather limited (see, for instance, [6, 7, 14–18]). This paucity of studies has arisen from the mathematical difficulty of obtaining an analytical solution to the wave equation of the stringent geometry being considered; in fact, analytical GF computations apply to spacetimes in which the metric components are independent of time. It is also worth noting that although BSs are defined as a higher-dimensional generalization of a BH, in which the event horizon is topologically equivalent to $S^2 \times S^1$ and spacetime is asymptotically $M^{D-1} \times S^1$, four-dimensional (4D) BSs are also derived. Lemos and Santos [19–21] showed that cylindrically symmetric static solutions, with a negative cosmological constant, of the Einstein–Maxwell equations admit charged 4D BSs. A rotating version of the charged 4D BSs [22, 23] exhibits features similar to the Kerr-Newman BH in spherical topology. The problem of analyzing GFs of scalar fields from charged 4D BSs has recently been discussed by Ahmed and Saifullah [24]. An interesting point about GFs has been reported by [25]: BH energy loss during Hawking radiation depends, crucially, on the GF and the particles' degrees of freedom.

As mentioned above, further study of the GFs of BSs is required. To fill this literature gap, in the current study, we considered dilatonic 5D BS [26], which is a solution to the Einstein–Yang–Mills–Born–Infeld–dilaton (EYMBID) theory. We analytically studied its GF, ACS, and DR for massive scalar fields; however, we considered tachyonic scalar particles instead of ordinary ones. The main reason for this consideration is that using ordinary mass in the

Klein–Gordon equation (KGE) of the dilatonic $5D$ BS (as will be explained in detail later) leads to the diverging of GFs. Roughly speaking, this is due to the flux of the propagating waves of the ordinary massive scalar fields. Namely, once the scalar fields to be considered belong to the massive ordinary particles, the incoming SI flux becomes zero. The latter remark implies that detectable radiation emitted from a dilatonic $5D$ BS spacetime belongs to the massive tachyonic scalar fields. Therefore, the current study focuses on the wave dynamics of tachyonic particles moving in dilatonic $5D$ BS spacetime. However, using tachyonic modes in $5D$ geometry should not be seen as nonphysical; instead they should be considered as the imaginary mass fields rather than faster-than-light particles [27]. First, Feinberg [28, 29] proposed that tachyonic particles could be quanta of a quantum field with imaginary mass. It was soon realized that excitations of *such imaginary mass fields do not in fact propagate faster than light* [30].

Following the idea of Kaluza–Klein [31], any $4D$ physical trajectory is the projection of higher-dimensional worldlines. Effective $4D$ worldlines associated with massive particles are causality constrained to be timelike. However, the corresponding higher-dimensional worldlines need not be exclusively timelike, which gives rise to a topological classification of physical objects. In particular, elementary particles in a $5D$ geometry should be viewed as tachyonic modes. The existence of tachyons in higher dimensions has been thoroughly studied by Davidson and Owen [32]. Furthermore, the reader may refer to [33] to understand tachyon condensation in the evaporation process of a BS. To find the analytical GF, ACS, and DR, we have shown how to obtain the complete analytical solution to the massive KGE in the geometry of a dilatonic $5D$ BS.

Our work is organized as follows. Following this introduction, a brief overview of the geometry of the dilatonic $5D$ BS is provided in Section 2. Section 3 describes the KGE of the tachyonic fields in the dilatonic $5D$ BS geometry; we present the exact solution of the radial equation in terms of hypergeometric functions. In Section 4, we compute the GF and consequently the ACS and DR of the dilatonic $5D$ BS, respectively. We then graphically exhibit the results of the ACS and DR. Section 5 concludes with the final remarks drawn from our study.

2. Dilatonic $5D$ BS in EYMBID Theory

$D\ (= d+1)$-dimensional action in the EYMBID theory is given by [26]

$$I_{EYMBID} = -\frac{1}{16\pi G_{(D)}}\int_{\mathscr{M}} d^d x\sqrt{-g} \cdot \left[\mathscr{R} - \frac{4(\nabla\psi)^2}{D-2} + 4\chi^2 e^{-b\psi}\left(1-\sqrt{1+\frac{Fe^{2b}}{2\chi^2}}\right)\right], \tag{1}$$

where ψ is the dilaton field, χ denotes the Born–Infeld parameter [34], and $b = -(4/(d-2))\alpha$ with the dilaton parameter $\alpha = 1/\sqrt{d-1}$. $G_{(D)}$ represents the D-dimensional Newtonian constant and its relation to its $4D$ form ($G_{(4)}$) is given by

$$G_{(D)} = G_{(4)}L^{D-4}, \tag{2}$$

where L is the upper limit of the compact coordinate ($\int_0^L dz = L$). Furthermore, $\mathscr{R}$ stands for the Ricci scalar and $F = F^{(\overline{a})}_{\lambda\rho}F^{(\overline{a})\lambda\rho}$ where the 2-form Yang–Mills field is given by

$$F^{(a)} = dA^{(\overline{a})} + \frac{1}{2\sigma}C^{(\overline{a})}_{(\overline{b})(\overline{c})}\left(A^{(\overline{b})}\wedge A^{(\overline{c})}\right), \tag{3}$$

with $C^{(\overline{a})}_{(\overline{b})(\overline{c})}$ and σ being structure and coupling constants, respectively. The Yang–Mills potential $A^{(a)}$ is defined by following the Wu-Yang ansatz [35]

$$A^{(\overline{a})} = \frac{Q}{r^2}\left(x_i dx_j - x_j dx_i\right), \tag{4}$$

$$r^2 = \sum_{i=1}^{d-1} x_i^2, \qquad 2 \le j+1 \le i \le d-1,\ 1 \le \overline{a} \le \frac{(d-1)(d-2)}{2}, \tag{5}$$

where Q is the Yang–Mills charge. The solution for the dilaton is as follows:

$$\psi = -\frac{(d-2)}{2}\frac{\alpha\ln r}{\alpha^2+1}. \tag{6}$$

On the other hand, the line-element of the dilatonic $5D$ BS is given by [26]

$$ds_0^2 = -\frac{f(r)}{\beta}d\widetilde{t}^2 + \frac{\beta dr^2}{rf(r)} + rd\widetilde{z}^2 + \beta\left(d\theta^2 + \sin^2\theta d\phi^2\right), \tag{7}$$

where $f(r) = r - r_+$ and $\beta = 4Q^2/3$. r_+ represents the outer event horizon having the following $(d+1)$–dimensional form:

$$\frac{32}{L^{d-4}}\left(\frac{Q^2 d}{d-1}\right)^{(d-2)/2} = r_+^{(d(d-2)+2)/d}. \tag{8}$$

Because, in our case, $D = 5$ (i.e., $d = 4$), the horizon becomes

$$r_+ = 4\beta^{2/5} = 4.488Q^{4/5}. \tag{9}$$

After rescaling the metric (7)

$$ds^2 = \frac{ds_0^2}{\beta} = -\frac{f(r)}{\beta^2}d\widetilde{t}^2 + \frac{dr^2}{rf(r)} + \frac{r}{\beta}d\widetilde{z}^2 + d\theta^2 + \sin^2\theta d\phi^2, \tag{10}$$

and in sequel assigning $\widetilde{t}$ and $\widetilde{z}$ coordinates to the new coordinates

$$\widetilde{t} \longrightarrow \beta t, \qquad \widetilde{z} \longrightarrow \beta z, \tag{11}$$

we get the metric that will be used in our computations:

$$ds^2 = -f(r)\,dt^2 + \frac{dr^2}{rf(r)} + \beta r dz^2 + d\theta^2 + \sin^2\theta d\phi^2. \quad (12)$$

It is worth noting that the surface gravity [36] of the dilatonic $5D$ BS can be evaluated by

$$\kappa^2 = -\frac{1}{2}\nabla^{\rho}\Upsilon^{\sigma}\nabla_{\rho}\Upsilon_{\sigma}\Big|_{r=r_+}, \quad (13)$$

in which Υ^{μ} represents the timelike Killing vector:

$$\Upsilon^{\mu} = [1, 0, 0, 0, 0]. \quad (14)$$

Then, (13) results in

$$\kappa = \frac{\sqrt{r}f'}{2}\Big|_{r=r_+} = \frac{\sqrt{r_+}}{2}, \quad (15)$$

where the prime denotes the derivative with respect to r. Furthermore, the associated Hawking temperature is expressed by

$$T_H = \frac{\kappa}{2\pi} = \frac{\sqrt{r_+}}{4\pi}. \quad (16)$$

It is important to remark that the Hawking temperature of the dilatonic BS given in (40) of [26] is incorrect. The authors of [26] computed the Hawking temperature of the dilatonic $5D$ BS considering the metric to be symmetric, which is not the case since $g_{tt} \neq 1/g_{rr}$. Meanwhile, it is obvious that dilatonic $5D$ BS (12) has a nonasymptotically flat structure. Therefore, it possesses a quasilocal mass [37–39], which can be computed as follows:

$$M_{QL} = \frac{1}{6}\beta r_h^{3/2} L = \frac{4}{3}\beta^{8/5} L \cong 2.113 Q^{16/5} L. \quad (17)$$

Thus, the first law of thermodynamics is satisfied:

$$dM_{QL} = T_H dS_{BH}, \quad (18)$$

where S_{BH} denotes the Bekenstein-Hawking entropy [36], which takes the following form for the dilatonic $5D$ BS:

$$\begin{aligned} S_{BH} &= \frac{A_H}{4} = \frac{1}{4}\beta r_h \int_0^{\pi} \sin\theta d\theta \int_0^{2\pi} d\phi \int_0^{L} dz \\ &= \pi\beta L r_h. \end{aligned} \quad (19)$$

3. Wave Equation of a Massive Scalar Tachyonic Field in Dilatonic $5D$ BS

As the scalar waves being studied belong to the massive scalar tachyons, the corresponding KGE is given by

$$\left[\Box - (i\mu)^2\right]\Psi(t, \mathbf{r}) = \left[\Box + \mu^2\right]\Psi(t, \mathbf{r}) = 0. \quad (20)$$

We chose the ansatz as follows:

$$\Psi(t, \mathbf{r}) = R(r)\, Y_l^m(\theta, \phi)\, e^{ikz} e^{-i\omega t}, \quad (21)$$

where $Y_l^m(\theta, \phi)$ is the usual spherical harmonics and k is a constant. After making straightforward calculations, we obtained the radial equation as follows:

$$rf\frac{\ddot{R}}{R} + \frac{\dot{R}}{R}\left(f + r\dot{f}\right) + \frac{\omega^2}{f} - \frac{k^2}{\beta r} + \mu^2 - \lambda = 0, \quad (22)$$

where $\lambda = l(l+1)$ and a dot mark denotes a derivative with respect to r. Multiplying each term by $r\beta f(r)R(r)$ and using the ansatz $y = (r_+ - r)/r_+$, which in turn implies $r = r_+ - yr_+$, one gets

$$\begin{aligned} &y(1-y)R'' + (1-2y)R' \\ &\quad + \left[\frac{\omega^2}{yr_+} + \frac{k^2}{\beta(1-y)r_+} - \mu^2 + \lambda\right]R = 0, \end{aligned} \quad (23)$$

where prime denotes derivative with respect to y. Setting

$$\left[\frac{\omega^2}{yr_+} + \frac{k^2}{\beta(1-y)r_+} - \mu^2 + \lambda\right] = \frac{A^2}{y} - \frac{B^2}{1-y} + C, \quad (24)$$

one can obtain

$$\begin{aligned} A &= -\frac{\omega}{2\kappa}, \\ B &= \frac{ik}{2\kappa\sqrt{\beta}}, \\ C &= \lambda - \mu^2. \end{aligned} \quad (25)$$

Equation (23) can be solved by comparing it with the standard hypergeometric differential equation [40] which admits the following solution:

$$\begin{aligned} R &= \xi_1(-y)^{iA}(1-y)^{-B} F(a, b; c; y) \\ &\quad + \xi_2(-y)^{-iA}(1-y)^{-B} F(\alpha, \varsigma; \eta; y), \end{aligned} \quad (26)$$

where

$$a = \frac{1}{2}\left(1 + \sqrt{1+4C}\right) + iA - B, \quad (27)$$

$$b = \frac{1}{2}\left(1 - \sqrt{1+4C}\right) + iA - B, \quad (28)$$

$$c = 1 + 2iA. \quad (29)$$

And

$$\alpha = a - c + 1, \quad (30)$$

$$\varsigma = b - c + 1, \quad (31)$$

$$\eta = 2 - c. \quad (32)$$

According to our calculations, to have a nonzero SI incoming flux (see (53)) or nondivergent GF (see (55)), $\sqrt{1+4C}$ must be imaginary. To this end, we must impose the following condition:

$$4\mu^2 > 4\lambda + 1, \quad (33)$$

such that

$$\sqrt{1+4C} = i\tau, \tag{34}$$

where

$$\tau = \sqrt{4\mu^2 - 4\lambda - 1}, \quad \tau \in \mathbb{R}. \tag{35}$$

To obtain a physically acceptable solution, we must terminate the outgoing solution at the horizon, which can be simply done by imposing $\xi_1 = 0$. Thus, the physical radial solution reduces to

$$R = \xi_2 (-y)^{-iA} (1-y)^{-B} F(\alpha, \varsigma; \gamma; y). \tag{36}$$

It should be noted that checking the forms of (36) both at the horizon and at SI is essential. Section 3 shows that both are needed for the evaluation of the GF. For the near horizon (NH) where $y \longrightarrow 0$, one can state

$$R_{NH} = \xi_2 (-y)^{-iA}, \tag{37}$$

which implies that the purely ingoing plane wave reads

$$\psi_{NH} = \xi_2 e^{-i\omega(\hat{r}_* + t)} e^{ikz}, \tag{38}$$

where

$$\begin{aligned} r^* &= \int \frac{dr}{\sqrt{r}f} \longrightarrow \\ \hat{r}_* &= \lim_{r\longrightarrow r_+} r^* \simeq \frac{\ln(-y)}{\sqrt{r_+}}, \\ \hat{r}_* &\simeq \frac{1}{2\kappa} \ln(-y) \Longrightarrow \\ y &= -e^{2\kappa\hat{r}_*}. \end{aligned} \tag{39}$$

On the other hand, for $y \longrightarrow \infty$, the inverse transformation of the hypergeometric function is given by [41]

$$\begin{aligned} F(\alpha, \varsigma; \eta; y) = (-y)^{-\alpha} & \frac{\Gamma(\eta)\Gamma(\varsigma-\alpha)}{\Gamma(\varsigma)\Gamma(\eta-\alpha)} \\ & \times F\left(\alpha, \alpha+1-\eta; \alpha+1-\varsigma; \frac{1}{y}\right) \\ & + (-y)^{-\varsigma} \frac{\Gamma(\eta)\Gamma(\alpha-\varsigma)}{\Gamma(\alpha)\Gamma(\eta-\varsigma)} \\ & \times F\left(\varsigma, \varsigma+1-\eta; \varsigma+1-\alpha; \frac{1}{y}\right), \end{aligned} \tag{40}$$

which yields the following asymptotic solution:

$$\begin{aligned} R_{SI} \simeq \xi_2 (-y)^{-iA-B-\alpha} & \frac{\Gamma(\eta)\Gamma(\varsigma-\alpha)}{\Gamma(\varsigma)\Gamma(\eta-\alpha)} \\ & + \xi_2 (-y)^{-iA-B-\varsigma} \frac{\Gamma(\eta)\Gamma(\alpha-\varsigma)}{\Gamma(\alpha)\Gamma(\eta-\varsigma)}. \end{aligned} \tag{41}$$

To express (41) in a more compact form, let us perform the following simplifications. Considering

$$-iA - B - \alpha = -\frac{1}{2}(1 + i\tau), \tag{42}$$

together with

$$-iA - B - \varsigma = -\frac{1}{2}(1 - i\tau), \tag{43}$$

and letting $x = -y$, the radial equation for $r \longrightarrow \infty$ takes the form

$$\begin{aligned} R_{SI} = \frac{1}{\sqrt{x}} \Bigg[& \xi_2 x^{-i\tau/2} \frac{\Gamma(\eta)\Gamma(\varsigma-\alpha)}{\Gamma(\varsigma)\Gamma(\eta-\alpha)} \\ & + \xi_2 x^{i\tau/2} \frac{\Gamma(\eta)\Gamma(\alpha-\varsigma)}{\Gamma(\alpha)\Gamma(\eta-\varsigma)} \Bigg]. \end{aligned} \tag{44}$$

One can express x in terms of the tortoise coordinate at SI as

$$\begin{aligned} r^* &= \int \frac{dr}{\sqrt{r}f} \longrightarrow \\ \hat{r}_* &= \lim_{r\longrightarrow\infty} r^* \simeq -\frac{2}{\sqrt{r}}, \end{aligned} \tag{45}$$

such that

$$\begin{aligned} x &= r - r_+ \\ x|_{r\longrightarrow\infty} &\simeq r = 4e^{-2\hat{r}_*}, \end{aligned} \tag{46}$$

where $\hat{r}_* = \ln r^*$. Therefore, we have

$$R_{SI} = \frac{1}{\sqrt{r}} \left[\Lambda_1 e^{i\hat{r}_*\tau} + \Lambda_2 e^{-i\hat{r}_*\tau} \right], \tag{47}$$

where

$$\Lambda_1 = 2^{-i\tau} \xi_2 \frac{\Gamma(\eta)\Gamma(\varsigma-\alpha)}{\Gamma(\varsigma)\Gamma(\eta-\alpha)}, \tag{48}$$

$$\Lambda_2 = 2^{-i\tau} \xi_2 \frac{\Gamma(\eta)\Gamma(\alpha-\varsigma)}{\Gamma(\alpha)\Gamma(\eta-\varsigma)}. \tag{49}$$

Thus, the asymptotic wave solution becomes

$$\psi_{SI} = \frac{e^{ikz}}{\sqrt{r}} \left[\Lambda_1 e^{i(\hat{r}_*\tau - \omega t)} + \Lambda_2 e^{-i(\hat{r}_*\tau + \omega t)} \right]. \tag{50}$$

4. Radiation of Dilatonic $5D$ BS

4.1. The Flux Computation. In this section, we compute the ingoing flux at the horizon ($r \longrightarrow r_+$) and the asymptotic flux for the SI region ($r \longrightarrow \infty$). The evaluation of these flux values will enable us to calculate the GF and, subsequently, the ACS and DR.

The NH-flux can be calculated via [42, 43]

$$\digamma_{NH} = \frac{A_{BH}}{2i} \left(\overline{\psi}_{NH} \partial r_* \psi_{NH} - \psi_{NH} \partial r_* \overline{\psi}_{NH} \right), \tag{51}$$

which, after a few manipulations, can be written as

$$\mathsf{F}_{NH} = -4\pi\beta\left|\xi_2\right|^2 r_+. \tag{52}$$

The incoming flux at SI is computed via

$$\mathsf{F}_{SI} = \frac{A_{BH}}{2i}\left(\overline{\psi}_{SI}\partial r_*\psi_{SI} - \psi_{SI}\partial r_*\overline{\psi}_{SI}\right). \tag{53}$$

Having performed the steps to evaluate the derivatives with respect to the tortoise coordinate, the incoming flux at SI takes the form

$$\mathsf{F}_{SI} = -4\pi\beta\left|\Lambda_2\right|^2 \tau. \tag{54}$$

It is important to remark that if we were dealing with the standard particles rather than tachyons, τ (35) would be imaginary, i.e., $\tau \longrightarrow i\tau$, and therefore SI incoming wave would lead this flux evaluation (53) to be zero. This would indicate the existence of a divergent GF.

4.2. ACS of Dilatonic 5D BS. The GF of the dilatonic 5D BS is obtained by the following expression [6, 8]:

$$\gamma^{l,k} = \frac{\mathsf{F}_{NH}}{\mathsf{F}_{SI}} = \frac{-4\pi\beta\left|\xi_2\right|^2 r_+}{-4\pi\beta\left|\Lambda_2\right|^2 \tau}, \tag{55}$$

which is nothing but

$$\gamma^{l,k} = \frac{\left|\xi_2\right|^2 r_+}{\left|\Lambda_2\right|^2 \tau}. \tag{56}$$

After a few manipulations, with (see [40])

$$\left|\Gamma(iy)\right|^2 = \frac{\pi}{y\sinh(\pi y)}, \tag{57}$$

$$\left|\Gamma(1+iy)\right|^2 = \frac{\pi y}{\sinh(\pi y)}, \tag{58}$$

$$\left|\Gamma\left(\frac{1}{2}+iy\right)\right|^2 = \frac{\pi}{\cosh(\pi y)}, \tag{59}$$

(56) can be presented as

$$\gamma^{l,k} = \frac{\kappa r_+}{\omega}\left(e^{2\pi\omega/\kappa} - 1\right)\Xi, \tag{60}$$

where

$$\Xi = \frac{e^{2\pi\tau} - 1}{\left[e^{\pi(\tau+\omega/\kappa - k/\kappa\sqrt{\beta})} + 1\right]\left[e^{\pi(\tau+\omega/\kappa + k/\kappa\sqrt{\beta})} + 1\right]}. \tag{61}$$

To evaluate the ACS of the dilatonic 5D BS concerned, we follow the study of [44]. Thus, one can get the ACS expression in 5D as follows:

$$\sigma^{l,k} = \frac{4\pi(l+1)^2}{\omega^3}\gamma^{l,k}, \tag{62}$$

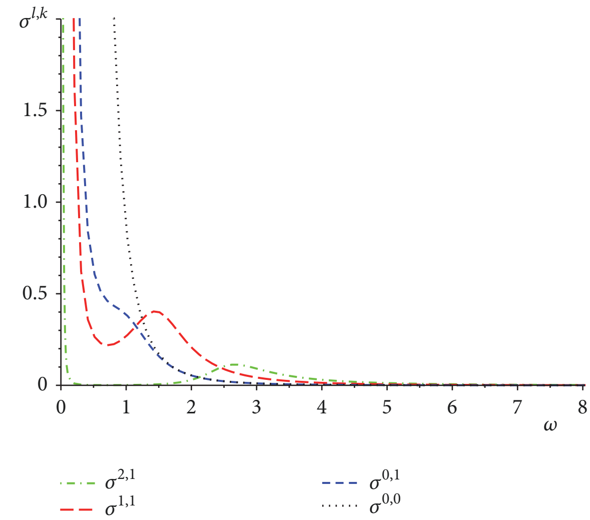

Figure 1: Plots of the ACS ($\sigma^{l,k}$) versus frequency ω. The plots are governed by (63). The configuration of the dilatonic 5D BS is as follows: $\mu = 3$ and $Q = 0.2$.

which, in our case, becomes

$$\sigma^{l,k} = \frac{4\pi(l+1)^2\kappa r_+}{\omega^4}\left(e^{2\pi\omega/\kappa} - 1\right)\Xi. \tag{63}$$

Furthermore, one can also get the total ACS as follows [45]:

$$\sigma_{abs}^{Total} = \sum_{l=0}^{\infty}\sigma^{l,k}. \tag{64}$$

In Figure 1, the relationship between absorption ACS and frequency is examined; the figure is drawn based on (63). In the high frequency regime, all ACSs tend to vanish by following the same curve. Unlike the high frequency regime, ACSs diverge in the low frequency regime as $\omega \longrightarrow 0$. As a final remark, negative $\sigma^{l,k}$ behavior has not been observed in our graphical analyses, which means that superradiance does not occur [46], as expected (as the dilatonic 5D BS (12) does not rotate).

4.3. DR of Dilatonic 5D BS. The final step follows from the ACS evaluation. The DR of the dilatonic 5D BS can be computed via [8]

$$\Gamma_{DR}^{l,k} = \frac{\sigma^{l,k}}{e^{2\pi\omega/\kappa} - 1} = \frac{4\pi(l+1)^2\kappa r_+}{\omega^4}\Xi. \tag{65}$$

Figure 2 shows how the DR behaves with respect to the frequency. By taking (65) as the reference, the plots for increasing l are illustrated. In the high frequency regime, all DRs fade in the same way. In the low frequency regime, DRs tend to diverge. However, it can be observed that when l has larger values, the corresponding DR diverges when ω is much closer to zero.

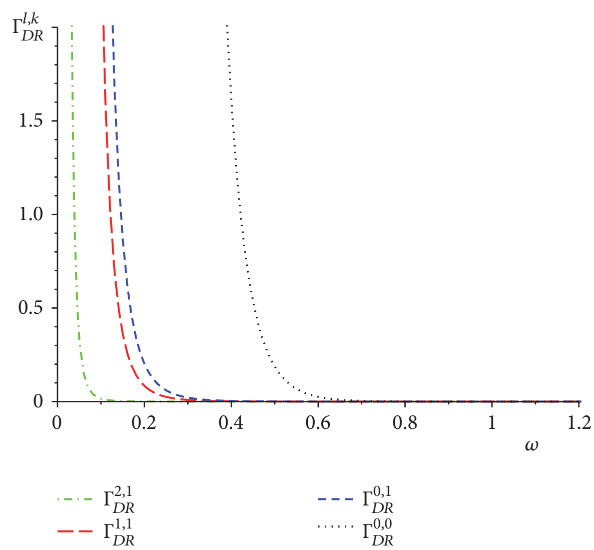

FIGURE 2: Plots of the DR ($\Gamma_{DR}^{l,k}$) versus frequency ω. The plots are governed by (65). The configuration of the dilatonic $5D$ BS is as follows: $\mu = 3$ and $Q = 0.2$.

5. Conclusion

This article evaluated the GF, ACS, and DR for the dilatonic $5D$ BS geometry arising from the EYMBID theory. As a result of the analytical method we followed, it was shown that the radiation of the dilatonic $5D$ BS spacetime can only be caused by tachyons. The crucial point here is that if standard scalar particles had been used rather than tachyonic ones, zero incoming flux at SI would have been obtained, which would lead to the diverging of the GF. Therefore, in a way, we were forced to use tachyons to solve this problem, and this carries great importance as it implies that the fifth dimension could be directly linked to tachyons. In short, according to our analytical method, we obtained results (compatible with boundary conditions) when the radiation of the dilatonic $5D$ BS was provided by the tachyons.

In future study, we want to extend our analysis to the Dirac equation for the geometry of the dilatonic $5D$ BS. Hence, we are planning to undertake similar analysis for fermions and compare the results with scalar ones.

Data Availability

The data that support the findings of this study are available from the corresponding author upon reasonable request.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

Acknowledgments

We wish to thank Prof. Dr. Mustafa Halilsoy for drawing our attention to this problem and for his helpful comments and suggestions. This work is supported by Eastern Mediterranean University through the project BAPC-04-18-01.

References

[1] S. W. Hawking, “Black hole explosions?” *Nature*, vol. 248, no. 5443, pp. 30-31, 1974.

[2] S. W. Hawking, “Particle creation by black holes,” *Communications in Mathematical Physics*, vol. 43, no. 3, pp. 199–220, 1975, erratum: ibid, vol. 46, p. 206, 1976.

[3] S. Creek, O. Efthimiou, P. Kanti, and K. Tamvakis, “Greybody factors for brane scalar fields in a rotating black hole background,” *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 75, no. 8, 084043, 17 pages, 2007.

[4] P. Boonserm, “Rigorous bounds on Transmission, Reflection, and Bogoliubov coefficients,” 2009, https://arxiv.org/abs/0907.0045.

[5] C. Hoyos and P. Koroteev, “On the null energy condition and causality in lifshitz holography,” *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 82, Article ID 109905, 2010.

[6] T. Harmark, J. Natario, and R. Schiappa, “Greybody Factors for d-Dimensional Black Holes,” *Advances in Theoretical and Mathematical Physics*, vol. 14, no. 3, pp. 727–793, 2010.

[7] R. Kallosh, A. Linde, and B. Vercnocke, “Natural inflation in supergravity and beyond,” *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 90, Article ID 041303(R), 2014.

[8] A. Sánchez-Bretaño, M. Callejo, M. Montero, Á. L. Alonso-Gómez, M. J. Delgado, and E. Isorna, “Performing a hepatic timing signal: glucocorticoids induce gper1a and gper1b expression and repress gclock1a and gbmal1a in the liver of goldfish,” *Journal of Comparative Physiology B: Biochemical, Systemic, and Environmental Physiology*, vol. 186, no. 1, pp. 73–82, 2016.

[9] G. Panotopoulos and A. Rincon, “Greybody factors for a minimally coupled massless scalar field in Einstein-Born-Infeld dilaton spacetime,” *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 96, no. 2, 025009, 7 pages, 2017.

[10] Y. Zhang, C. Li, and B. Chen, “Greybody factors for a spherically symmetric Einstein-Gauss-Bonnet–de Sitter black hole,” *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 97, Article ID 044013, 2018.

[11] J. Kunz, “Black holes in higher dimensions (black strings and black rings),” in *Proceedings of the MG13 Meeting on General Relativity*, pp. 568–581, Stockholm University, Sweden, 2015.

[12] R. Gregory and R. Laflamme, “The instability of charged black strings and p-branes,” *Nuclear Physics. B. Theoretical, Phenomenological, and Experimental High Energy Physics. Quantum Field Theory and Statistical Systems*, vol. 428, no. 1-2, pp. 399–434, 1994.

[13] R. Gregory and R. Laflamme, “Black strings and p-branes are unstable,” *Physical Review Letters*, vol. 70, no. 19, pp. 2837–2840, 1993.

[14] I. R. Klebanov and S. D. Mathur, “Black hole greybody factors and absorption of scalars by effective strings,” *Nuclear Physics. B. Theoretical, Phenomenological, and Experimental High Energy Physics. Quantum Field Theory and Statistical Systems*, vol. 500, no. 1-3, pp. 115–132, 1997.

[15] J. Callan, S. S. Gubser, I. R. Klebanov, and A. A. Tseytlin, “Absorption of fixed scalars and the D-brane approach to black holes,” *Nuclear Physics. B. Theoretical, Phenomenological, and Experimental High Energy Physics. Quantum Field Theory and Statistical Systems*, vol. 489, no. 1-2, pp. 65–94, 1997.

[16] M. Cvetic and F. Larsen, "General rotating black holes in string theory: greybody factors and event horizons," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 56, no. 8, pp. 4994–5007, 1997.

[17] I. R. Klebanov and M. Krasnitz, "Fixed scalar greybody factors in five and four dimensions," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 55, no. 6, pp. R3250–R3254, 1997.

[18] S. S. Gubser and I. R. Klebanov, "Emission of charged particles from four- and five-dimensional black holes," *Nuclear Physics. B. Theoretical, Phenomenological, and Experimental High Energy Physics. Quantum Field Theory and Statistical Systems*, vol. 482, no. 1-2, pp. 173–186, 1996.

[19] J. P. Lemos, "Two-dimensional black holes and planar general relativity," *Classical and Quantum Gravity*, vol. 12, no. 4, pp. 1081–1086, 1995.

[20] J. P. S. Lemos, "Three dimensional black holes and cylindrical general relativity," *Physics Letters B*, vol. 353, p. 46, 1995.

[21] N. O. Santos, "Solution of the vacuum Einstein equations with nonzero cosmological constant for a stationary cylindrically symmetric spacetime," *Classical and Quantum Gravity*, vol. 10, no. 11, pp. 2401–2406, 1993.

[22] J. P. S. Lemos and V. T. Zanchin, "Rotating charged black strings and three-dimensional black holes," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 54, no. 6, pp. 3840–3853, 1996.

[23] R. G. Cai and Y. Z. Zhang, "Black plane solutions in four-dimensional spacetimes," *Physical Review D*, vol. 54, p. 4891, 1996.

[24] J. Ahmed and K. Saifullah, "Greybody factor of scalar fields from black strings," *The European Physical Journal C*, vol. 77, p. 885, 2017.

[25] E. K. Akhmedov and J. Pulido, "Solar neutrino oscillations and bounds on neutrino magnetic moment and solar magnetic field," *Physics Letters B*, vol. 553, no. 1-2, pp. 7–17, 2003.

[26] S. H. Mazharimousavia and M. Halilsoy, "Non-Abelian magnetic black strings versus black holes," *The European Physical Journal Plus*, vol. 131, p. 138, 2016.

[27] A. Sen, "Rolling Tachyon," *Journal of High Energy Physics*, vol. 2002, no. 04, p. 048, 2002.

[28] G. Feinberg, "Possibility of faster-than-light particles," *Physical Review A: Atomic, Molecular and Optical Physics*, vol. 159, no. 5, pp. 1089–1105, 1967.

[29] L. Wolfenstein, "Neutrino oscillations in matter," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 17, no. 9, pp. 2369–2374, 1978.

[30] Y. Aharonov, A. Komar, and L. Susskind, "Superluminal Behavior, Causality, and Instability," *Physical Review A: Atomic, Molecular and Optical Physics*, vol. 182, no. 5, pp. 1400–1403, 1969.

[31] O. Klein, "Quantentheorie und fünfdimensionale relativitätstheorie," *Zeitschrift für Physik*, vol. 37, no. 12, pp. 895–906, 1926.

[32] A. Davidson and D. A. Owen, "Elementary particles as higher-dimensional tachyons," *Physics Letters B*, vol. 177, no. 1, pp. 77–81, 1986.

[33] G. Horowitz, "Tachyon condensation and black strings," *Journal of High Energy Physics*, vol. 2005, no. 08, p. 091, 2005.

[34] M. Born and L. Infeld, "Foundations of the new field theory," *Proceedings of the Royal Society A Mathematical, Physical and Engineering Sciences*, vol. 144, no. 852, pp. 425–451, 1934.

[35] T. T. Wu and C. N. Yang, *in Properties of Matter Under Unusual Conditions*, New York, NY, USA, Interscience, 1969, edited by H. Mark and S. Fernbach.

[36] R. M. Wald, *General Relativity*, University of Chicago Press, London, UK, 1984.

[37] J. D. Brown and J. W. York Jr., "Quasilocal energy and conserved charges derived from the gravitational action," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 47, no. 4, pp. 1407–1419, 1993.

[38] S. Bose and N. Dadhich, "Brown-York quasilocal energy, gravitational charge, and black hole horizons," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 60, no. 6, Article ID 064010, 7 pages, 1999.

[39] S. Chakraborty and N. Dadhich, "Brown-York quasilocal energy in Lanczos-Lovelock gravity and black hole horizon," *Journal of High Energy Physics*, vol. 12, p. 003, 2015.

[40] M. Abramowitz and I. A. Stegun, *Handbook of Mathematical Functions, with Formulas, Graphs, and Mathematical Tables*, Dover, New York, NY, USA, 1972.

[41] S. Y. Slavyanov and W. Lay, *Special Functions, A Unified Theory Based on Singularities*, Oxford Mathematical Monographs, Oxford, UK, 2000.

[42] S. Fernando, "Greybody factors of charged dilaton black holes in 2 + 1 dimensions," *General Relativity and Gravitation*, vol. 37, no. 3, pp. 461–481, 2005.

[43] M. Andrews, M. Lewandowski, M. Trodden, and D. Wesley, "Distinguishing k-defects from their canonical twins," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 82, Article ID 105006, 2010.

[44] P. Kanti, J. Grain, and A. Barrau, "Bulk and brane decay of a-dimensional Schwarzschild-de SITter black hole: scalar radiation," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 71, no. 10, 104002, 18 pages, 2005.

[45] W. G. Unruh, "Absorption cross section of small black holes," *Physical Review D: Particles, Fields, Gravitation and Cosmology*, vol. 14, no. 12, p. 3251.

[46] S. Chandrasekhar, *The Mathematical Theory of Black Holes*, The Clarendon Press, New York, NY, USA, 1983.