

# Hawking Radiation of Photons in a Vaidya–de Sitter Black Hole

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Hawking evaporation of photons in a Vaidya–de Sitter black hole is investigated by using the method of generalized tortoise coordinate transformation. Both the location and the temperature of the event horizon depend on the time. It is shown that Hawking radiation of photons exists only for the complex Maxwell scalar  $\phi_0$  in the advanced Eddington–Finkelstein coordinate system. This asymmetry of Hawking radiation for different components of Maxwell fields probably arises from the asymmetry of space-time in the advanced Eddington–Finkelstein coordinate system. It is shown that the black body radiant spectrum of photons resembles that of Klein–Gordon particles.

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**KEY WORDS:** Hawking radiation; Maxwell equation; Vaidya-type black hole; generalized tortoise coordinate transformation; Newman–Penrose formalism; null tetrad.

## 1. INTRODUCTION

Hawking's investigation of quantum effects (Hawking, 1974, 1975) interpreted as the emission of a thermal spectrum of particles by a black hole event horizon sets a significant landmark in black hole physics. In the last few decades, much work has been done on the Hawking evaporation of black holes in some spherically symmetric and nonstatic spacetimes (Kim *et al.*, 1989; Li and Zhao, 1993; Ma and Yang, 1993, 1997; Yang and Zhao, 1994; Zhang *et al.*, 1996; Zhao *et al.*, 1992; Zhu *et al.*, 1994a,b). In a recent paper (Wu and Cai, 2001) (here refer to Paper I), we reexamined the Hawking effect of Dirac particles in a Vaidya-type black hole by means of the generalized tortoise transformation method. We considered simultaneously the asymptotic behaviors of the first-order and second-order forms of Dirac equations near the event horizon, and eliminated the crossing terms of the first-order derivatives in the second-order equations by using the relations between the first-order derivatives of the radial Dirac equations. We showed that the Hawking radiation exists only for  $P_2$  and  $Q_1$  components of Dirac spinors in

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virtue of the restriction imposed by the limiting form of its first-order equations. We conceived that this asymmetry of the Hawking radiation for different spinorial components probably originates from the asymmetry of spacetime in the advanced Eddington–Finkelstein coordinate system.

In this paper, we deal with the thermal radiation of photons in a Vaidya–de Sitter spacetime. The method used here is the same as that presented in Paper I, namely, we consider the asymptotic behaviors of the first-order and second-order forms of Maxwell equations in the vicinity of the event horizon, and recast each second-order equation to a standard wave equation near the event horizon. The location and the temperature of the event horizon are shown to be dependent on the time. The black body radiation spectrum of photons resembles that of Klein–Gordon scalar particles. We find that because of the restrictions put on the Hawking radiation by the limiting form of the first-order Maxwell equations, not all Maxwell complex scalars but  $\phi_0$  displays the property of thermal radiation. This asymmetry of Hawking radiation for different field components in a Vaidya-type spacetime is thought to be a common feature shared by all particles with higher spins.

The paper is organized as follows: In Section 2, the explicit form of sourceless Maxwell equations within the framework of Newman–Penrose formalism (Newman and Penrose, 1962) are written out by choosing an appropriate null tetrad in the Vaidya–de Sitter geometry. By using the method of generalized tortoise coordinate transformation, the event horizon equation is extracted in Section 3 from the asymptotic forms of the radial parts of the first-order Maxwell equations near the event horizon. Then the second-order radial equations are manipulated in Section 4 by the same procedure and recast into a standard wave equation near the event horizon; in the meanwhile, an exact expression of the “surface gravity” of the event horizon is also obtained by adjusting the parameter  $\kappa$ . Section 5 is devoted to deriving the thermal radiation spectrum of photons from the event horizon. Finally we present some discussions in Section 6.

## 2. SOURCELESS MAXWELL EQUATIONS

The metric of a Vaidya black hole with a cosmological constant  $\Lambda$  is given in the advanced Eddington–Finkelstein coordinate system by

$$ds^2 = 2 dv(G dv - dr) - r^2(d\theta^2 + \sin^2 \theta d\varphi^2), \quad (1)$$

where  $2G = 1 - 2M/r - \Lambda r^4/3$ , in which mass  $M(v)$  of the hole is a function of the advanced time  $v$ .

The geometry of this Petrov type-D spacetime is characterized by two kinds of surfaces of particular interest: the apparent horizon  $r_{\text{AH}} = 2M$  (coincides with the timelike limit surfaces  $r_{\text{TLS}}$ ) and the event horizon  $r_{\text{EH}} = r_{\text{H}}$ . The apparent horizon is the outer most trapped surface, while the event horizon is necessarily a null surface  $r = r(v)$  that satisfies the null-surface conditions  $g^{ij} \partial_i F \partial_j F = 0$

and  $F(v, r) = 0$ . Traditionally the latter is calculated approximately by the simple physical condition that the photons at the event horizon are stuck or unaccelerated in the sense that  $\dot{r} \approx 0$ . An effective method for determining the location of the event horizon of a dynamic black hole is the generalized tortoise coordinate transformation (GTCT), with which we can apply it to the null-hypersurface equation  $g^{ij} \partial_i F \partial_j F = 0$  and then take the limits approaching the event horizon. However in Paper I, the event horizon equation is extracted from the asymptotic forms of the radial parts of the first-order Dirac equations near the event horizon. In this paper, we derive it by the same procedure but with the Maxwell equations in place of the Dirac equations here.

When the back reaction of the massless spin-1 test particles on the background geometry is neglected, the electromagnetic field equation is given by the Maxwell equation on a fixed spacetime (1). In order to write out its explicit form in the Newman–Penrose (Newman and Penrose, 1962) formalism, we establish the following complex null-tetrad system  $\{l, n, m, \bar{m}\}$  that satisfies the orthogonal conditions  $l \cdot n = -m \cdot \bar{m} = 1$ . Thus the covariant one-forms can be written as

$$\begin{aligned}
 l &= dv, & m &= \frac{-r}{\sqrt{2}}(d\theta + i \sin \theta d\varphi), \\
 n &= G dv - dr, & \bar{m} &= \frac{-r}{\sqrt{2}}(d\theta - i \sin \theta d\varphi),
 \end{aligned}
 \tag{2}$$

and their corresponding directional derivatives are

$$\begin{aligned}
 D &= -\frac{\partial}{\partial r}, & \delta &= \frac{1}{\sqrt{2}r} \left( \frac{\partial}{\partial \theta} + \frac{i}{\sin \theta} \frac{\partial}{\partial \varphi} \right), \\
 \Delta &= \frac{\partial}{\partial v} + G \frac{\partial}{\partial r}, & \bar{\delta} &= \frac{1}{\sqrt{2}r} \left( \frac{\partial}{\partial \theta} - \frac{i}{\sin \theta} \frac{\partial}{\partial \varphi} \right).
 \end{aligned}
 \tag{3}$$

It is not difficult to determine the nonvanishing Newman–Penrose complex spin coefficients (Newman and Penrose, 1962) in the above null tetrad as

$$\begin{aligned}
 \rho &= 1/r, & \gamma &= -G_{,r}/2 = -dG/2 dr, \\
 \mu &= G/r, & \beta &= -\alpha = \cot \theta / (2\sqrt{2}r).
 \end{aligned}
 \tag{4}$$

Inserting for the appropriate spin coefficients into the sourceless Maxwell equations (Carmeli, 1982; Chandrasekhar, 1983; Newman and Penrose, 1962) in the Newman–Penrose formalism (Newman and Penrose, 1962)

$$\begin{aligned}
 (D - 2\rho)\phi_1 - (\bar{\delta} + \tilde{\pi} - 2\alpha)\phi_0 &= -\tilde{\kappa}\phi_2, \\
 (\delta - 2\tau)\phi_1 - (\Delta + \mu - 2\gamma)\phi_0 &= -\sigma\phi_2, \\
 (D + 2\epsilon - \rho)\phi_2 - (\bar{\delta} + 2\tilde{\pi})\phi_1 &= -\tilde{\lambda}\phi_0, \\
 (\delta + 2\beta - \tau)\phi_2 - (\Delta + 2\mu)\phi_1 &= -\tilde{\nu}\phi_0,
 \end{aligned}
 \tag{5}$$

we obtain

$$\begin{aligned}
 \left(\frac{\partial}{\partial r} + \frac{2}{r}\right)\phi_1 + \frac{1}{\sqrt{2}r}\mathcal{L}_1\phi_0 &= 0, \\
 \left(\mathcal{D} + \frac{G}{r} + G_{,r}\right)\phi_0 - \frac{1}{\sqrt{2}r}\mathcal{L}_0^\dagger\phi_1 &= 0, \\
 \left(\frac{\partial}{\partial r} + \frac{1}{r}\right)\phi_2 + \frac{1}{\sqrt{2}\rho}\mathcal{L}_0\phi_1 &= 0, \\
 \left(\mathcal{D} + \frac{2G}{r}\right)\phi_1 - \frac{1}{\sqrt{2}\rho^*}\mathcal{L}_1^\dagger\phi_2 &= 0.
 \end{aligned} \tag{6}$$

Here we have defined operators:

$$\begin{aligned}
 \mathcal{D} &= \frac{\partial}{\partial v} + G \frac{\partial}{\partial r}, \\
 \mathcal{L}_n &= \frac{\partial}{\partial \theta} + n \cot \theta - \frac{i}{\sin \theta} \frac{\partial}{\partial \varphi}, \\
 \mathcal{L}_n^\dagger &= \frac{\partial}{\partial \theta} + n \cot \theta + \frac{i}{\sin \theta} \frac{\partial}{\partial \varphi}.
 \end{aligned}$$

By substituting  $\Phi_0 = r\phi_0$ ,  $\Phi_1 = \sqrt{2}r^2\phi_1$ , and  $\Phi_2 = r\phi_2$  into Eq. (6), we have

$$\begin{aligned}
 \frac{\partial}{\partial r}\Phi_1 + \mathcal{L}_1\Phi_0 &= 0, & 2r^2(\mathcal{D} + G_{,r})\Phi_0 - \mathcal{L}_0^\dagger\Phi_1 &= 0, \\
 2r^2\frac{\partial}{\partial r}\Phi_2 + \mathcal{L}_0\Phi_1 &= 0, & \mathcal{D}\Phi_1 - \mathcal{L}_1^\dagger\Phi_2 &= 0.
 \end{aligned} \tag{7}$$

### 3. EVENT HORIZON

Equation (7) can be decoupled as

$$\Phi_0 = R_0(v, r)S_0(\theta, \varphi), \quad \Phi_1 = R_1(v, r)S_1(\theta, \varphi), \quad \Phi_2 = R_2(v, r)S_2(\theta, \varphi)$$

to the radial part

$$\begin{aligned}
 \frac{\partial}{\partial r}R_1 + \lambda R_0 &= 0, & 2r^2(\mathcal{D} + G_{,r})R_0 + \lambda R_1 &= 0, \\
 2r^2\frac{\partial}{\partial r}R_2 + \lambda R_1 &= 0, & \mathcal{D}R_1 + \lambda R_2 &= 0,
 \end{aligned} \tag{8}$$

and the angular part

$$\begin{aligned}
 \mathcal{L}_1S_0 &= \lambda S_1, & \mathcal{L}_0^\dagger S_1 &= -\lambda S_0, \\
 \mathcal{L}_0S_1 &= \lambda S_2, & \mathcal{L}_1^\dagger S_2 &= -\lambda S_1,
 \end{aligned} \tag{9}$$

where  $\lambda = \sqrt{\ell(\ell + 1)}$  is a separation constant. All functions  $S_0(\theta, \varphi)$ ,  $S_1(\theta, \varphi)$ , and  $S_2(\theta, \varphi)$  are, respectively, spin-weighted spherical harmonics,  ${}_p Y_{\ell m}(\theta, \varphi)$  with spin weight  $p = 1, 0, -1$ , satisfying the following equation (Goldberg *et al.*, 1968)

$$\left[ \frac{\partial^2}{\partial \theta^2} + \cot \theta \frac{\partial}{\partial \theta} + \frac{1}{\sin^2 \theta} \frac{\partial^2}{\partial \varphi^2} + \frac{2ip \cos \theta}{\sin^2 \theta} \frac{\partial}{\partial \varphi} - p^2 \cot^2 \theta + p + (\ell - p)(\ell + p + 1) \right] {}_p Y_{\ell m}(\theta, \varphi) = 0. \quad (10)$$

As to the Hawking radiation, we should be concerned about the asymptotic behaviors of the radial part of Eq. (8) in the vicinity of the event horizon. Because the Vaidya–de Sitter spacetime is spherically symmetric, we can introduce as a working ansatz the following GTCT as in the Paper I

$$r_* = r + \frac{1}{2\kappa} \ln[r - r_H(v)], \quad v_* = v - v_0, \quad (11)$$

where  $r_H = r(v)$  is the location of the event horizon, and  $\kappa$  is an adjustable parameter and is unchanged under the tortoise transformation. The parameter  $v_0$  is an arbitrary constant that characterizes the initial instant of the hole. From formula (11), we can deduce some useful relations for the derivatives as follows:

$$\frac{\partial}{\partial r} = \frac{\partial}{\partial r_*} + \frac{1}{2\kappa(r - r_H)} \frac{\partial}{\partial r_*}, \quad \frac{\partial}{\partial v_*} = \frac{\partial}{\partial v_*} - \frac{\dot{r}_H}{2\kappa(r - r_H)} \frac{\partial}{\partial r_*}.$$

The quantities  $\dot{r}_H = \partial r_H / \partial v$  is the rate of the event horizon varying in time  $v$ . It describes the evolution of the black hole event horizon in the time, which reflects the presence of quantum ergosphere near the event horizon.

Now let us consider first the asymptotic behaviors of Eq. (8) near the event horizon. Under the transformations (11), Eq. (8) can be reduced to the following forms

$$\begin{aligned} \frac{\partial}{\partial r_*} R_1 &= 0, & 2r_H^2(\dot{r}_H - G) \frac{\partial}{\partial r_*} R_0 &= 0, \\ 2r_H^2 \frac{\partial}{\partial r_*} R_2 &= 0, & (\dot{r}_H - G) \frac{\partial}{\partial r_*} R_1 &= 0, \end{aligned} \quad (12)$$

after having taken the  $r \rightarrow r_H(v_0)$  and  $v \rightarrow v_0$  limits.

From Eq. (12), we know that  $R_1(r_*)$  and  $R_2(r_*)$  are regular on the event horizon,

$$\frac{\partial}{\partial r_*} R_1 = \frac{\partial}{\partial r_*} R_2 = 0. \quad (13)$$

Thus, a reasonable solution to Eq. (12) is that the derivatives  $\frac{\partial}{\partial r_*} R_0$  do not vanish. The sole possibility we are left with for the existence of a nontrivial solution of  $R_0$

is (as for  $r_H \neq 0$ )

$$2G(r_H) - 2\dot{r}_H = 0, \tag{14}$$

which determines the location of the event horizon. It is interesting to note that the event horizon equation (14) coincides with that inferred from the null-surface equation  $g^{ij} \partial_i F \partial_j F = 0$ . Because  $r_H$  depends on time  $v$ , the location of the event horizon and the shape of the black hole change with time.

#### 4. HAWKING TEMPERATURE

In the preceding section, we have deduced the event horizon equation from the limiting form of the separated radial part of the first-order Maxwell equations. Applying a similar procedure to its second-order forms, we can derive the Hawking temperature and the thermal radiation spectrum. A straightforward calculation gives the second-order radial equations

$$2r^2 \left[ G \frac{\partial^2}{\partial r^2} + \frac{\partial^2}{\partial v \partial r} + 2 \left( G_{,r} + \frac{G}{r} \right) \frac{\partial}{\partial r} + \frac{2}{r} \frac{\partial}{\partial v} + \frac{2G_{,r}}{r} + G_{,rr} \right] \times R_1 - \lambda^2 R_1 = 0, \tag{15}$$

$$2r^2 \left( G \frac{\partial^2}{\partial r^2} + \frac{\partial^2}{\partial v \partial r} + G_{,r} \frac{\partial}{\partial r} \right) R_1 - \lambda^2 R_1 = 0, \tag{16}$$

$$2r^2 \left( G \frac{\partial^2}{\partial r^2} + \frac{\partial^2}{\partial v \partial r} + \frac{2G}{r} \frac{\partial}{\partial r} \right) R_2 - \lambda^2 R_2 = 0. \tag{17}$$

Given the GTCT in Eq. (11) and after some tedious calculations, Eqs. (15)–(17) have the following limiting forms near the event horizon  $r = r_H$

$$\left\{ \left[ \frac{A}{2\kappa} + 4G(r_H) - 2\dot{r}_H \right] \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} + \left[ -A + 4G_{,r}(r_H) + \frac{4G(r_H) - 4\dot{r}_H}{r_H} \right] \frac{\partial}{\partial r_*} \right\} R_0 = \left\{ \left( \frac{A}{2\kappa} - 2\dot{r}_H \right) \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} + [-A + 4G_{,r}(r_H)] \frac{\partial}{\partial r_*} \right\} R_0 = 0, \tag{18}$$

$$\left\{ \left[ \frac{A}{2\kappa} + 4G(r_H) - 2\dot{r}_H \right] \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} + [-A + 2G_{,r}(r_H)] \frac{\partial}{\partial r_*} \right\} R_1 = \left[ \left( \frac{A}{2\kappa} - 2\dot{r}_H \right) \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} \right] R_1 = 0, \tag{19}$$

$$\left\{ \left[ \frac{A}{2\kappa} + 4G(r_H) - 2\dot{r}_H \right] \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} + \left[ -A + \frac{4G(r_H)}{r_H} \right] \frac{\partial}{\partial r_*} \right\} R_2 = \left[ \left( \frac{A}{2\kappa} - 2\dot{r}_H \right) \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} \right] R_2 = 0, \quad (20)$$

when  $r$  approaches  $r_H(v_0)$  and  $v$  nears  $v_0$ . In the above calculations, we have used relations  $2G(r_H) = 2\dot{r}_H$  and  $\frac{\partial}{\partial r_*} R_1 = \frac{\partial}{\partial r_*} R_2 = 0$ . The coefficient  $A$  is an infinite form of  $0/0$  type with a finite result, treated by using the L' Hôpital rule,

$$A = \lim_{r \rightarrow r_H(v_0)} \frac{2G - 2\dot{r}_H}{r - r_H} = 2G_{,r}(r_H).$$

To recast Eqs. (18), (19), and (20) into a standard wave equation near the event horizon, we select the adjustable parameter  $\kappa$  in them such that it satisfies

$$\frac{A}{2\kappa} + 2G(r_H) = \frac{G_{,r}(r_H)}{\kappa} + 2\dot{r}_H \equiv 1, \quad (21)$$

which means the “surface gravity” of the horizon is

$$\kappa = \frac{G_{,r}(r_H)}{1 - 2G(r_H)} = \frac{G_{,r}(r_H)}{1 - 2\dot{r}_H}, \quad (22)$$

where we have used Eq. (14).

With such a parameter adjustment, these wave equations (18)–(20) can be reduced to

$$\left( \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} + 2C \frac{\partial}{\partial r_*} \right) R_0 = 0, \quad (23)$$

$$\left( \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} \right) R_1 = 0, \quad (24)$$

$$\left( \frac{\partial^2}{\partial r_*^2} + 2 \frac{\partial^2}{\partial r_* \partial v_*} \right) R_2 = 0, \quad (25)$$

where  $C = G_{,r}(r_H)$ . Equations (23)–(25) have a standard form of wave equation in the vicinity of the event horizon. We point out that the above parameter adjustment is an important step in our discussions.

## 5. THERMAL RADIATION SPECTRUM

Combining Eqs. (24) and (25) with  $\frac{\partial}{\partial r_*} R_1 = \frac{\partial}{\partial r_*} R_2 = 0$ , we know that  $R_1$  and  $R_2$  are independent of  $r_*$  near the event horizon. The solutions  $R_1 \sim e^{-i\omega v_*}$  and  $R_2 \sim e^{-i\omega v_*}$  indicate that Hawking radiation does not exist for  $\Phi_1$  and  $\Phi_2$ .

Now separating variables from Eq. (23) as  $R_0 = R_0(r_*)e^{-i\omega v_*}$ , one gets

$$R_0'' = 2(i\omega - C)R_0', \quad R_0 = c_1 e^{2(i\omega - C)r_*} + c_2. \quad (26)$$

The incoming wave solution and the outgoing wave solution to Eq. (23) are, respectively,

$$\begin{aligned} R_0^{\text{in}} &\sim e^{-i\omega v_*}, \\ R_0^{\text{out}} &\sim e^{-i\omega v_*} e^{2(i\omega-C)r_*}, \quad r > r_{\text{H}}. \end{aligned} \tag{27}$$

Near the event horizon, we have  $r_* \sim (1/2\kappa) \ln(r - r_{\text{H}})$ . Clearly, the outgoing wave  $R_0^{\text{out}}(r > r_{\text{H}})$  is not analytic at the event horizon  $r = r_{\text{H}}$ , but can be analytically extended from the exterior of the hole into the interior of the hole by the lower complex  $r$  plane

$$(r - r_{\text{H}}) \rightarrow (r_{\text{H}} - r) e^{-i\pi}$$

to

$$\widetilde{R}_0^{\text{out}} = e^{-i\omega v_*} e^{2(i\omega-C)r_*} e^{i\pi C/\kappa} e^{\pi\omega/\kappa}, \quad r < r_{\text{H}}. \tag{28}$$

Following the method of Damour–Ruffini–Sannan’s (Damour and Ruffini, 1976; Sannan, 1988), the relative scattering probability of the outgoing wave at the event horizon and the thermal radiation spectrum of photons from the event horizon of the black hole are easily obtained

$$\left| \frac{R_0^{\text{out}}}{\widetilde{R}_0^{\text{out}}} \right|^2 = e^{-2\pi\omega/\kappa}, \quad \langle \mathcal{N}_\omega \rangle \sim \frac{1}{e^{\omega/T} - 1}, \tag{29}$$

in which  $m$  is the azimuthal quantum number, and the Hawking temperature is

$$T = \frac{\kappa}{2\pi} = \frac{1}{4\pi r_{\text{H}}} \cdot \frac{Mr_{\text{H}} - \Lambda r_{\text{H}}^4/3}{Mr_{\text{H}} - \Lambda r_{\text{H}}^4/6}. \tag{30}$$

The thermal radiation spectrum (29) due to the Bose–Einstein statistics of photons shows that the black hole emits radiation just like a black body emitting scalar particles. The temperature depends on the time and is consistent with that derived from the investigation of the thermal radiation of Dirac particles in a Vaidya–de Sitter black hole (Wu and Cai, 2001) with a vanishing electric charge ( $Q = 0$ ).

## 6. CONCLUSIONS

In this paper, we have studied the Hawking radiation of photons in a Vaidya–de Sitter black hole. The location and the temperature of the event horizon given by Eqs. (14) and (22), respectively, depend on the advanced time  $v$ . They can recover the well-known results previously derived in the discussion on the Hawking evaporation of Klein–Gordon and Dirac particles in the same spacetime. Equation (29) shows that the thermal radiation spectra of photons have the same form as that of Klein–Gordon particles in a Vaidya-type black hole with a cosmological constant  $\Lambda$ .



In summary, we have dealt with the asymptotic behaviors of the separated radial equations near the event horizon, not only its first-order form but also its second-order form. We find that the limiting form of its first-order equations puts very strong restriction on the Hawking effect, that is, not all components of Maxwell complex scalars but  $\phi_0$  displays the property of thermal radiation. As is revealed in Paper I, we argued that this asymmetry of Hawking radiation for different components of Maxwell fields probably stem from the asymmetry of spacetimes in the advanced Eddington–Finkelstein coordinate. We think this is a common character shared by the thermal radiation of particles with higher spins in any Vaidya-type spherically symmetric black hole.

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