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A Jet Appearing When a Black Hole Event Horizon Touches the Rindler Horizon

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The Hawking effect of a nonuniformly rectilinearly accelerating Kinnersley black hole is studied. Its horizons are rotationally symmetric. Its Hawking temperature depends not only on the time, but also on the polar angle. When a Kinnersley black hole touches its Rindler horizon, the Hawking temperature at the contact point is reduced to zero. But at the opposite pole of the black hole, the Hawking temperature increases rapidly. A jet appears as a tail of the accelerating black hole.

1. INTRODUCTION

Recently, we suggested a new method to determine the location and the temperature of event horizons of nonstationary black holes (Zhao and Dai, 1992; Zhu *et al.*, 1995; Zhao and Li, 1993; Zhao *et al.*, 1994). By means of the new method, we have given those for some spherically symmetric nonstatic black holes. The results are consistent with those obtained by calculating the vacuum expectation values of the renormalized energy-momentum tensors (Hiscock, 1986; Balbinot, 1986). Furthermore, the new method is more exact and more convenient than the old one.

Now we want to deal with the Hawking effect of an axially symmetric nonstationary black hole—nonuniformly rectilinearly accelerating Kinnersley (1969) black hole. It is impossible to get the Hawking temperature of non-spherically symmetric, nonstationary black holes using the calculation of the energy-momentum tensors, but it is possible and easy by means of the new method. In addition, we want to know what happens when the black hole

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horizon touches the Rindler horizon in the Kinnersley space-time (Unruh, 1976).

In Section 2 we study the Hawking effect in Kinnersley space-time, and show that the horizons are rotationally symmetric and that the Hawking temperature depends not only on the time, but also on the polar angles. In Section 3 we study the Hawking–Unruh effect of the Rindler horizon when the mass of the Kinnersley black hole vanishes. Section 4 contains the very interesting result that a jet will appear as a tail of the accelerating black hole when the black hole touches its Rindler horizon. In Section 5 we give conclusion and discussion.

2. HAWKING EFFECT IN KINNERSLEY SPACE-TIME

The line element of the nonuniformly rectilinearly accelerating Kinnersley black hole space-time is (Kinnersley, 1969)

$$ds^2 = (1 - 2ar\cos\theta - r^2f^2 - 2mr^{-1})\,dv^2 - 2\,dv\,dr - 2r^2f\,dv\,d\theta - r^2\,d\theta^2 - r^2\sin^2\theta\,d\phi^2 \tag{1}$$

where $f = -a(v)\sin\theta$.

Both acceleration a and mass m depend on the advanced Eddington time v. The north pole $\theta = 0$ points toward the direction of acceleration. By means of the null-surface condition

$$g^{\mu\nu}\left(\frac{\partial F}{\partial x^\mu}\right)\left(\frac{\partial F}{\partial x^\nu}\right) = 0 \tag{2}$$

we have the event horizon equation (Zhao *et al.*, 1993)

$$2\dot{r}_H - \left(1 - 2ar_H\cos\theta - \frac{2m}{r_H}\right) - 2fr'_H - \frac{r'^2_H}{r^2_H} = 0 \tag{3}$$

r_H is the location of the event horizon, $\dot{r}_H = (\partial r/\partial v)_H$, and $r'_H = (\partial r/\partial\theta)_H$.

The Klein–Gordon equation can be given as

$$-2\frac{\partial^2\Phi}{\partial v\,\partial r} - \frac{2}{r}\frac{\partial\Phi}{\partial v} - \left(1 - 2ar\cos\theta - \frac{2m}{r}\right)\frac{\partial^2\Phi}{\partial r^2} - \frac{1}{r^2}(2r + 6ar^2\cos\theta - 2m)\frac{\partial\Phi}{\partial r} + 2f\frac{\partial^2\Phi}{\partial r\,\partial\theta} + \frac{2f}{r}\frac{\partial\Phi}{\partial\theta} - 2a\cos\theta\frac{\partial\Phi}{\partial r} - \frac{1}{r^2}\frac{\partial^2\Phi}{\partial\theta^2} - \frac{\operatorname{ctg}\theta}{r^2}\frac{\partial\Phi}{\partial\theta} - \frac{1}{r^2\sin^2\theta}\frac{\partial^2\Phi}{\partial\phi^2} - \mu^2\Phi = 0 \tag{4}$$

With the generalized tortoise coordinate transformation (Zhao and Li, 1992; Zhu *et al.*, 1995; Zhao and Li, 1993; Zhao *et al.*, 1994)

$$r_* = r + \frac{1}{2\kappa}\ln[r - r_H(v, \theta)], \qquad v_* = v - v_0, \qquad \theta_* = \theta - \theta_0 \quad (5)$$

equation (4) can be written as

$$\frac{\{2\dot{r}_H - (1 - 2ar\cos\theta - 2mr^{-1})[2\kappa(r - r_H) + 1] - 2fr'_H\} \\ r^2[2\kappa(r - r_H) + 1] - r'^2_H}{2\kappa(r - r_H)[2\kappa(r - r_H) + 1]r^2}$$

$$\times \frac{\partial^2\Phi}{\partial r_*^2} - 2\frac{\partial^2\Phi}{\partial r_*\partial v_*} + \left\{\frac{2r'_H}{r^2[2\kappa(r - r_H) + 1]} + 2f\right\}\frac{\partial^2\Phi}{\partial r_*\partial\theta_*}$$

$$+ \left[1 + \frac{1}{2\kappa(r - r_H)}\right]^{-1}$$

$$\times \left\{\frac{-2\dot{r}_H}{2\kappa(r - r_H)^2} + \frac{2\dot{r}_H}{2\kappa r(r - r_H)} + \frac{1 - 2ar\cos\theta - 2mr^{-1}}{2\kappa(r - r_H)^2}\right.$$

$$- \frac{2(r - 8ar^2\cos\theta - m)}{r^2}\left[1 + \frac{1}{2\kappa(r - r_H)}\right] + \frac{2fr'_H}{2\kappa(r - r_H)^2}$$

$$- \frac{2fr'_H}{2\kappa r(r - r_H)}$$

$$\left. + \frac{r'^2_H + r''_H(r - r_H)}{2\kappa r^2(r - r_H)^2} + \frac{r'_H\cos\theta}{2\kappa r^2(r - r_H)\sin\theta}\right\}\frac{\partial\Phi}{\partial r_*}$$

$$- \left[1 + \frac{1}{2\kappa(r - r_H)}\right]^{-1}$$

$$\times \left\{\frac{2}{r}\frac{\partial\Phi}{\partial v_*} + \frac{1}{r}\left(-2f + \frac{\operatorname{ctg}\theta}{r}\right)\frac{\partial\Phi}{\partial\theta_*} + \frac{1}{r^2}\frac{\partial^2\Phi}{\partial\theta_*^2}\right.$$

$$\left. + \frac{1}{r^2\sin^2\theta}\frac{\partial^2\Phi}{\partial\phi^2} + \mu^2\Phi\right\} = 0 \quad (6)$$

where κ is an adjustable parameter in the tortoise transformation. It is the surface gravity of the event horizon when the space-time is stationary.

When $r \to r_H(v_0, \theta_0)$, $v \to v_0$, and $\theta \to \theta_0$, the equation can be reduced to

$$\alpha\frac{\partial^2\Phi}{\partial r_*^2} + 2\frac{\partial^2\Phi}{\partial r_*\partial\theta_*} + B\frac{\partial^2\Phi}{\partial r_*\partial\theta_*} - G\frac{\partial\Phi}{\partial r_*} = 0 \quad (7)$$

where

$$\alpha \equiv \lim_{\substack{r\to r_H(v_0,\theta_0)\\ v\to v_0\\ \theta\to\theta_0}} \frac{\{2\dot{r}_H - (1 - 2ar\cos\theta - 2mr^{-1})[2\kappa(r - r_H) + 1] - 2fr'_H\} \times r^2[2\kappa(r - r_H) + 1] - r'^2_H}{-2\kappa(r - r_H)[2\kappa(r - r_H) + 1]r^2} \tag{8}$$

$$B = -2\left(f + \frac{r'_H}{r_H^2}\right)_{\substack{v\to v_0\\ \theta\to\theta_0}} \tag{9}$$

$$G = \left(-\frac{2}{r_H} + 16\, a\cos\theta + \frac{2m}{r_H^2} - \frac{r'^2_H}{r_H^3} + \frac{r'_H}{r_H^2}\,\mathrm{ctg}\,\theta\right)_{\substack{v\to v_0\\ \theta\to\theta_0}} \tag{10}$$

Here, we have used equation (3). We select the adjustable parameter κ as

$$\kappa = \frac{1}{2r_H}\, \frac{m/r_H^2 - a\cos\theta - r'^2_H/r_H^3}{m/r_H^2 + a\cos\theta + r'^2_H/2r_H^3}\bigg|_{\substack{v\to v_0\\ \theta\to\theta_0}} \tag{11}$$

Then we have $\alpha = 1$, and equation (7) can be reduced to

$$\frac{\partial^2\Phi}{\partial r_*^2} + 2\frac{\partial^2\Phi}{\partial r_*\partial v_*} + B\frac{\partial^2\Phi}{\partial r_*\partial\theta_*} - G\frac{\partial\Phi}{\partial r_*} = 0 \tag{12}$$

Separating variables as

$$\Phi = R(r_*)\Theta(\theta_*)\exp[-i\omega v_* + in\phi] \tag{13}$$

we can verify that the radial wave solutions of equation (12) are, respectively,

$$\phi_{\mathrm{in}} = \exp[-i\omega v_*] \tag{14}$$

$$\phi_{\mathrm{out}} = \exp[-i\omega v_* + Gr_* + 2i\omega r_*] \tag{15}$$

ϕ_{in} is the ingoing wave, while ϕ_{out} is the outgoing wave. Near the event horizon r_H, ϕ_{out} can be rewritten as

$$\phi_{\mathrm{out}} = \exp[-i\omega v_*](r - r_H)^{G/2\kappa}(r - r_H)^{i\omega/\kappa} \tag{16}$$

It is not analytical at the horizon. By analytical continuation rotating $-\pi$ through the lower-half complex r-plane, we can extend ϕ_{out} from the outside of the black hole into the inside of the black hole (Damour, 1976)

$$\begin{cases} (r - r_H) \to |r - r_H|e^{-i\pi} = (r_H - r)e^{-i\pi} \\ \psi_{\mathrm{out}} \to \psi'_{\mathrm{out}} = e^{-i\omega v_* + Gr_* + 2i\omega r_*}e^{-i\pi G/2\kappa}e^{\pi\omega/\kappa} \end{cases} \tag{17}$$

The relative scattering probability of the outgoing wave at the horizon is

$$\left|\frac{\psi_{out}}{\psi'_{out}}\right|^2 = e^{-2\pi\omega/\kappa} \tag{18}$$

Then the spectrum of the Hawking radiation is (Sannan, 1988)

$$N_\omega = (e^{\omega/K_B T} \pm 1)^{-1} \tag{19}$$

where

$$T = \frac{\kappa}{2\pi K_B} = \frac{1}{2\pi K_B}\frac{1}{2r_H}\frac{m/r_H^2 - a\cos\theta - r_H'^2/r_H^3}{m/r_H^2 + a\cos\theta + r_H'^2/2r_H^3} \tag{20}$$

T is the Hawking temperature and K_B is Boltzmann's constant.

We see that both the location and the temperature of the event horizons in the nonuniformly rectilinearly accelerating Kinnersley space-time depend not only on the time, but also on the polar angle.

3. RINDLER HORIZON

When $m = 0$ but $a(v) \neq 0$, we give the Rindler horizon equation of a nonuniformly rectilinearly accelerating observer,

$$2\dot{r}_H - (1 - 2ar_H\cos\theta) - 2fr_H' - \frac{r_H'^2}{r_H^2} = 0 \tag{21}$$

where r_H is the location of the Rindler horizon. The Klein–Gordon equation can be given as

$$\begin{aligned} -2\frac{\partial^2\Phi}{\partial v\,\partial r} &- \frac{2}{r}\frac{\partial\Phi}{\partial v} - (1 - 2ar\cos\theta)\frac{\partial^2\Phi}{\partial r^2} \\ &- \frac{1}{r^2}(2r + 6ar^2\cos\theta)\frac{\partial\Phi}{\partial r} + 2f\frac{\partial^2\Phi}{\partial r\,\partial\theta} + \frac{2f}{r}\frac{\partial\Phi}{\partial\theta} - 2a\cos\theta\frac{\partial\Phi}{\partial r} \\ &- \frac{1}{r^2}\frac{\partial^2\Phi}{\partial\theta^2} - \frac{\text{ctg}\,\theta}{r^2}\frac{\partial\Phi}{\partial\theta} - \frac{1}{r^2\sin^2\theta}\frac{\partial^2\Phi}{\partial\phi^2} - \mu^2\Phi = 0 \end{aligned} \tag{22}$$

We deal with equation (22) outside the Rindler horizon. Because the region outside the Rindler horizon is $r < r_H$, the generalized tortoise coordinate transformation should be rewritten as

$$r_* = r + \frac{1}{2\kappa}\ln[r_H(v,\theta) - r], \qquad v_* = v - v_0, \qquad \theta_* = \theta - \theta_0 \tag{23}$$

so equation (22) can be written

$$\frac{\{2\dot{r}_H - (1 - 2ar\cos\theta)[2\kappa(r - r_H) + 1] - 2fr'_H\} \times r^2[2\kappa(r - r_H) + 1] - r'^2_H}{2\kappa(r - r_H)[2\kappa(r - r_H) + 1]r^2}$$
$$\times \frac{\partial^2\Phi}{\partial r_*^2} - 2\frac{\partial^2\Phi}{\partial r_*\partial v_*} + \left\{\frac{2r'_H}{r^2[2\kappa(r - r_H) + 1]} + 2f\right\}\frac{\partial^2\Phi}{\partial r_*\partial\theta_*}$$
$$+ \left[1 + \frac{1}{2\kappa(r - r_H)}\right]^{-1}$$
$$\times \left\{\frac{-2\dot{r}_H}{2\kappa(r - r_H)^2} + \frac{2\dot{r}_H}{2\kappa r(r - r_H)} + \frac{1 - 2ar\cos\theta}{2\kappa(r - r_H)^2}\right.$$
$$- \frac{2(r - 8ar^2\cos\theta)}{r^2}\left[1 + \frac{1}{2\kappa(r - r_H)}\right]$$
$$+ \frac{2fr'_H}{2\kappa(r - r_H)^2} - \frac{2fr'_H}{2\kappa r(r - r_H)} + \frac{r'^2_H + r''_H(r - r_H)}{2\kappa r^2(r - r_H)^2}$$
$$\left. + \frac{r'_H\cos\theta}{2\kappa r^2(r - r_H)\sin\theta}\right\}\frac{\partial\Phi}{\partial r_*}$$
$$- \left[1 + \frac{1}{2\kappa(r - r_H)}\right]^{-1} \cdot \left\{\frac{2}{r}\frac{\partial\Phi}{\partial v_*} + \frac{1}{r}\left(-2f\frac{\operatorname{ctg}\theta}{r}\right)\frac{\partial\Phi}{\partial\theta_*}\right.$$
$$\left. + \frac{1}{r^2}\frac{\partial^2\Phi}{\partial\theta_*^2} + \frac{1}{r^2\sin^2\theta}\frac{\partial^2\Phi}{\partial\phi^2} + \mu^2\Phi\right\} = 0 \tag{24}$$

When $r \to r_H(v_0, \theta_0)$, $v \to v_0$ and $\theta \to \theta_0$, the equation can be reduced to

$$\alpha\frac{\partial^2\Phi}{\partial r_*^2} + 2\frac{\partial^2\Phi}{\partial r_*\,\partial v_*} + B\frac{\partial^2\Phi}{\partial r_*\,\partial\theta_*} - G\frac{\partial\Phi}{\partial r_*} = 0 \tag{25}$$

where

$$\alpha \equiv \lim_{\substack{r\to r_H(v_0,\theta_0)\\ v\to v_0\\ \theta\to\theta_0}} \times \frac{\{2\dot{r}_H - (1 - 2ar\cos\theta)[2\kappa(r - r_H) + 1] - 2fr'_H\}r^2[2\kappa(r - r_H) + 1] - r'^2_H}{-2\kappa(r - r_H)[2\kappa(r - r_H) + 1]r^2} \tag{26}$$

$$B = -2\left(r + \frac{r'_H}{r_H^2}\right)_{\substack{v \to v_0 \\ \theta \to \theta_0}} \tag{27}$$

$$G = \left(-\frac{2}{r_H} + 16a \cos\theta - \frac{r'^2_H}{r_H^3} + \frac{r'_H}{r_H^2} \operatorname{ctg}\theta\right)_{\substack{v \to v_0 \\ \theta \to \theta_0}} \tag{28}$$

Here, we have used equation (21). We select the adjustable parameter κ as

$$\kappa = \frac{1}{2r_H} \frac{-a\cos\theta - r'^2_H/r_H^3}{a\cos\theta + r'^2_H/2r_H^3}\bigg|_{\substack{v \to v_0 \\ \theta \to \theta_0}} \tag{29}$$

Then we have $\alpha = 1$, and equation (25) can be reduced to

$$\frac{\partial^2\Phi}{\partial r_*^2} + 2\frac{\partial^2\Phi}{\partial r_* \partial v_*} + B\frac{\partial^2\Phi}{\partial r_* \partial \theta_*} - G\frac{\partial\Phi}{\partial r_*} = 0 \tag{30}$$

Separating variables as

$$\Phi = R(r_*)\Theta(\theta_*) \exp[-i\omega v_* + in\phi] \tag{31}$$

we can verify that the radial wave solutions of equation (30) are, respectively,

$$\phi_{\text{in}} = \exp[-i\omega v_*] \tag{32}$$

$$\phi_{\text{out}} = \exp[-i\omega v_* + Gr_* + 2i\omega r_*] \tag{33}$$

ϕ_{in} is the ingoing wave, while ϕ_{out} is the outgoing wave. Near the Rindler horizon r_H, ϕ_{out} can be rewritten as

$$\phi_{\text{out}} = \exp[-i\omega v_*](r_H - r)^{G/2\kappa}(r_H - r)^{i\omega/\kappa} \tag{34}$$

It is not analytical at the horizon. By analytical continuation rotating $+\pi$ through the lower-half complex r-plane, we can extend ϕ_{out} from the outside of the Rindler horizon, $r < r_H$, into its inside, where $r > r_H$,

$$\begin{cases} (r_H - r) \to |r_H - r|e^{i\pi} = (r - r_H)e^{i\pi} \\ \psi_{\text{out}} \to \psi'_{\text{out}} = e^{i\omega v_* + Gr_* + 2i\omega r_*} e^{i\pi G/2\kappa} e^{-\pi\omega/\kappa} \end{cases} \tag{35}$$

The relative scattering probability of the outgoing wave at the horizon is

$$\left|\frac{\psi_{\text{out}}}{\psi'_{\text{out}}}\right|^2 = e^{2\pi\omega/\kappa} \tag{36}$$

Then the spectrum of the Hawking radiation is

$$N_\omega = (e^{\omega/K_B T} \pm 1)^{-1} \tag{37}$$

where

$$T = \frac{-\kappa}{2\pi K_B} = \frac{1}{2\pi K_B}\frac{1}{2r_H}\frac{a\cos\theta + r_H'^2/r_H^3}{a\cos\theta + r_H'^2/2r_H^3} \tag{38}$$

T is the Hawking–Unruh temperature, K_B is Boltzmann's constant.

We see that both the location and the temperature of the Rindler horizons in the nonuniformly rectilinearly accelerating Rindler space-time depend not only on the time, but also on the polar angle.

When $m = 0$ and $a = \text{const}$ we obtain the Rindler horizon of a uniformly rectilinearly accelerating observer,

$$1 - 2ar_H\cos\theta + 2fr_H' + \frac{r_H'^2}{r_H^2} = 0 \tag{39}$$

It is a paraboloid of revolution

$$r_H = \frac{1}{a(1 + \cos\theta)} \tag{40}$$

Its Hawking temperature is a constant, $\kappa = -a$,

$$T = \frac{a}{2\pi K_B} \tag{41}$$

When $m = 0$ and $a = a(v)$, the Rindler horizon will deviate from a paraboloid, and its Hawking temperature will depend on time v and the polar angle θ. When $a = 0$ but $m \neq 0$, from equation (3), we get the Vaidya black hole (Zhao and Dai, 1992; Balbinot, 1986)

$$r_H = \frac{2m}{1 - 2\dot{r}_H} \tag{42}$$

whose temperature depends on time v, but not on θ,

$$\kappa = \frac{1 - 2\dot{r}_H}{4m} \tag{43}$$

When $m = m(v) \neq 0$ and $a = a(v) \neq 0$, there exist two event horizons, the Rindler horizon and the black hole horizon. From equations (3) and (20), we know that these horizons are still rotationally symmetric and their Hawking–Unruh temperatures depend on both v and θ.

4. CONTACT OF BLACK HOLE AND RINDLER HORIZONS

Now let us consider what will happen when the black hole touches its Rindler horizon. First, we look for the polar angles, where r_H gets its maximum or minimum. With $r_H' = 0$, equation (3) can be reduced to

$$(2a \cos \theta_1) r_H^2 - (1 - 2\dot{r}_H) r_H + 2m = 0 \tag{44}$$

where θ_1 is the polar angle where r_H gets its extreme value. The solutions of equation (44) are

$$r_H = \frac{(1 - 2\dot{r}_H) \pm [(1 - 2\dot{r}_H)^2 - 16ma \cos \theta_1]^{1/2}}{4a \cos \theta_1} \tag{45}$$

They belong to the black hole horizon and the Rindler horizon, respectively. When $\dot{r}_H = 0$ and $ma << 1$, the above equation can be reduced to

$$r_{H1} \approx \frac{1}{2a \cos \theta_1}, \qquad r_{H2} \approx 2m \tag{46}$$

Apparently, r_{H2} belongs to the black hole horizon. Comparing equation (40) with equation (46), we know that r_{H1} belongs to the Rindler horizon, and $\theta_1 = 0$. Now, we show the extreme values of r_H,

$$r_{H1} = \frac{(1 - 2\dot{r}_{H1}) + \sqrt{(1 - 2\dot{r}_{H1})^2 - 16ma}}{4a}$$

$$r_{H2} = \frac{(1 - 2\dot{r}_{H2}) - \sqrt{(1 - 2\dot{r}_{H2})^2 - 16ma}}{4a} \tag{47}$$

When the black hole touches its Rindler horizon, i.e., $r_{H1} = r_{H2}$, we have

$$(1 - 2\dot{r}_H) = 16ma, \qquad a = \frac{m}{r_{H1}^2} = \frac{m}{r_{H2}^2} \tag{48}$$

and

$$r_{H1} = r_{H2} = \frac{1 - 2\dot{r}_H}{4a} = \frac{4m}{1 - 2\dot{r}_H} \tag{49}$$

Equation (20) is reduced to

$$T_2 = \frac{\kappa_2}{2\pi K_B} = \frac{1}{2\pi K_B} \frac{1}{2r_{H2}} \frac{m/r_{H2}^2 - a}{m/r_{H2}^2 + a} \tag{50}$$

and equation (38) is reduced to

$$T_1 = \frac{-\kappa_1}{2\pi K_B} = \frac{1}{2\pi K_B} \frac{1}{2r_{H1}} \frac{m/r_{H1}^2 - a}{m/r_{H1}^2 + a} \tag{51}$$

We see that the temperature of the contact point ($\theta = 0$, $r_H = 4m/(1 - 2\dot{r}_H)$) of the two horizons will be reduced to zero (Zhao *et al.*, 1994). This will violate the third law of thermodynamics.

On the other hand, at the opposite pole (south pole, $\theta = \pi$) of the black hole, the Hawking temperature will increase rapidly,

$$T_2 = \frac{\kappa_2}{2\pi K_B} = \frac{1}{2\pi K_B} \frac{1}{2r_H} \frac{m/r_H^2 + a - r_H'^2/r_H^3}{m/r_H^2 - a + r_H'^2/2r_H^3} \tag{52}$$

This shows that a thermal jet will appear as the tail of the accelerating black hole.

5. CONCLUSION AND DISCUSSION

The shape of the event horizon and the Rindler horizon of a nonuniformly rectilinearly accelerating Kinnersley black hole is rotationally symmetric and depends on both the time and the acceleration. Its Hawking temperature not only depends on the time, but also on the polar angle. When a Kinnersley black hole touches its Rindler horizon, the Hawking temperature at the contact point will be reduced to zero, but at the opposite pole of the black hole, the Hawking temperature will increase rapidly. A jet will appear as the tail of the accelerating black hole.

The "collision" between the black hole horizon and the Rindler horizon is very similar to the collision between two black holes. We infer that the Hawking temperature at the contact point of the two black holes will be reduced to zero, and that thermal jets will appear as tails of the two collision black holes. Every black hole will have a thermal jet tail.

REFERENCES

Balbinot, R. (1986). *Physical Review D*, **33**, 1611.

Damour, T., and Ruffini, R. (1976). *Physical Review D*, **14**, 332.

Hiscock, W. A. (1981). *Physical Review D*, **23**, 2813, 2823.

Kinnersley, W. (1969). *Physical Review*, **186**, 1335.

Sannan, S. (1988). *General Relativity and Gravitation*, **20**, 239.

Unruh, W. G. (1976). *Physical Review D*, **14**, 870.

Zhao Zheng and Dai Xianxin (1992). *Modern Physics Letters A*, **7**, 1771.

Zhao Zheng and Li Zhong-heng (1993). *Nuovo Cimento B*, **108**, 785.

Zhao Zheng, Yang Cheng-quan, and Ren Qin-an (1994). *General Relativity and Gravitation*, **26**, 1055.

Zhu Jian-yang, Bao Ai-dong, and Zhao Zheng (1995). *International Journal of Theoretical Physics*, **34**, 2049.