An Intuitive Understanding of Black-Hole Evaporation by Viewing it in Terms of More Familiar Quantum-Field Effects in Flat Spacetime

R. M. Nugayev
Kazan State Teachers Institute, Kazan, 420021, USSR

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The article is aimed at an intuitive understanding of the recently explored deep connections between thermal physics, quantum field theory and general relativity. The physical effects involved in particle creation by a black hole are viewed in terms of more familiar quantum-field effects in flat spacetime. Black hole evaporation is investigated in terms of temperature correction to the Casimir effect. The application of the Casimir effect results and those for accelerated mirrors reveals that a black hole should produce the blackbody radiation at a temperature that exactly coincides with Hawking's result. Its blackbody nature is due to the interaction of virtual positive-energy particles with the surface of a "cavity" formed by the Schwarzchild gravitational field potential barrier. The virtual particles are "squeezed out" by the contraction of the potential barrier and appear to an observer at J^+ as the real blackbody ones.

In previous papers ([2]–[4]) a programme of the reduction of particle creation by a black hole to quantum-field effects in flat space-time was initiated. The programme is based on the fact (R. H. Price) that the gravitational field of a black hole creates an effective potential barrier penetrable for high-frequency waves and impenetrable for waves with low frequency. The barrier is so well-localized near $r = 1.5$ Rg ($\mathrm{Rg} = 2\,GM/c^2$) that for the study of wave propagation we can consider the regions quite near the horizon and far away from it as "flat". Almost all the scattering takes place in the narrow region near $r = 1.5$ Rg. The consideration of the barrier peak as the surface of a reflecting sphere permitted to apply to a black hole the results of numerous Casimir-effect calculations. It appeared that the flow of negative Casimir energy should cause the hole to produce thermal radiation at a temperature that exactly coincides with the result of Hawking [1]. But the sequence of models designed to fit the evaporation process with increasing precision was unable to describe the very mechanism of particle creation. The reasons are obvious: all of them ignore the thickness of the potential barrier. In [2]–[4] the barrier was approximated by a thin shell. However, Fabbri [5] demonstrated that there exist two branches of turning points for a non-rotating black hole potential barrier:

$$r_1 = \frac{2}{\omega}\left[\frac{l(l+1)}{3}\right]^{1/2} \mathrm{Cos}\,\frac{\eta}{3},$$

$$r_2 = \frac{2}{\omega}\left[\frac{l(l+1)}{3}\right]^{1/2} \mathrm{Cos}\,\frac{\eta - 2\pi}{3},$$

where $\eta = \arccos\{-3\,\omega M[3/l(l+1)]^{1/2}\}$ and arccos denotes the principal value of the inverse trigonometric function, so that $2M \leqq r_1 \leqq r_2$. For instance, each (w, l) partial low-frequency wave has two turning points

$$r_1 = 2M\left[1 + \frac{4\,\omega^2 M^2}{l(l+1)} + O\left(\frac{\omega^4 M^4}{l^4}\right)\right],$$

$$r_2 = \frac{[l(l+1)]^{1/2}}{\omega} - M\left[1 + O\left(\frac{\omega M}{l}\right)\right], \qquad (1)$$

where $O(x)$ denotes a quantity of order x.

Consequently, for the purpose of investigating the interaction of virtual particles with the potential barrier surface it should be represented by two conducting concentric spheres. One of them is situated near the horizon while the other is far away from it. Each shell is an ideal conductor. So, the aim of the paper is to make the next step in describing the *evaporation mechanism* by means of an ideal model that is more realistic than the preceeding ones.

Reprint requests to Dr. Nugayev, Kazan State Teachers Institute, Kazan 420021, USSR.

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A) The results concerning temperature corrections to the Casimir effect (M. Fiertz, J. Mehra, R. Hargreaves, Brown, and Maclay et al.) were generalized by Tadaki and Takagi [6]. They considered two parallel infinite plane boundaries in the four-dimensional Minkowski spacetime. This system has two special directions (t, z) because of the presence of the boundaries and the heat bath. According to the symmetry of the system, the conservation law, and the tracelessness, the vacuum stress-tensor between the plates has the form

$$\langle T_{\mu\nu}\rangle = A\,\mathrm{diag}(-1,1,1,-3) + B\,\mathrm{diag}(3,1,1,1) + C\,\mathrm{diag}(1,0,0,1) + F(z)\,\mathrm{diag}(2,1,1,0)\,. \quad (2)$$

The first term in (2) represents the zero-temperature term and the second the Stephan-Boltzmann term. For a conformally coupled massless scalar field one has

$$A = \frac{\pi^2 \hbar c}{1440\, d^4}\,, \quad B = \frac{\pi^2 k^4 T^4}{90\, \hbar^3 c^3}$$

with d being the distance between the plates and T the temperature of each plate. The detailed form of $C = C(T)$ and $F = F(z)$ is given in [6].

In the low-temperature limit ($Td \ll 1$) the Stephan-Boltzmann term is negligible. The temperature correction is exponentially small, because the basis modes have an energy gap.

In the high-temperature limit ($Td \gg 1$) $\langle T_{\mu\nu}\rangle$ is dominated by the Stephan-Boltzmann value everywhere not close to the boundary. The behaviour near the boundary may be seen by considering the single boundary problem. In the limit $d \to \infty$ the result is

$$A = 0\,, \quad B = \frac{\pi^2 k^4 T^4}{90\, \hbar^3 c^3}\,, \quad C = 0\,. \quad (3)$$

The thermal average deviates from the Stephan-Boltzmann value near the boundary ($Z \ll I$) due to the T^4 term of $F(z)$:

$$F(z) = -\frac{\pi^2 T^4 k^4}{90\, \hbar^3 c^3} \cdot \frac{4}{3}\left\{1 - \frac{2}{7} z^2 + \frac{2}{35} z^4 + O(z^6)\right\}\,. \quad (4)$$

The calculations for the electromagnetic field are almost the same.

B) Levin, Polevoy, and Rytov ([7]) obtained, by means of the generalized Kirchhoff law, general expressions for the spectral and total Poynting vector for a fluctuation of the electromagnetic field in a plane vacuum gap between two arbitrary infinite media with different temperatures (for simplicity the media are assumed to be isotropic and spatially local). The Poynting vector is

$$P = \int_0^\infty p(\omega)\, \mathrm{d}\omega = \frac{1}{\pi^2} \int_0^\infty (\Pi_1 - \Pi_2)\, M\, \mathrm{d}\omega\,, \quad (5)$$

where

$$\Pi_i = \frac{\hbar\omega}{\exp(\hbar\omega/kT_i) - 1}\,, \quad i = 1, 2\,.$$

In vacuum ($\varepsilon_1 = \varepsilon_2 = \mu_1 = \mu_2 = k = 1$) for infinite separation ($d \to \infty$) one gets

$$M(\infty) = \omega^2/8c^2\,. \quad (6)$$

For

$$d = 0:\ M(0) = \omega^2/4c^2\,,$$

$$P(0) = \sigma_{\mathrm{SB}}(T_1^4 - T_2^4)\,, \quad (7)$$

where $T_1 > T_2$.

Thus, though each conductor is in thermal equilibrium with radiation, but each at its own temperature, the whole system is in a nonequilibrium state. Under these conditions *a flow of the fluctuating electromagnetic field from T_1 to T_2 ($T_1 > T_2$) dominates inside the cavity over the flow from T_2 to T_1.*

The authors of [7] point out that the generalized Kirchhoff's law contains an expression for the average oscillator energy $Q(w, T)$. Nevertheless, zero oscillations do not contribute the energy flow: $\Pi_i = Q(\omega, T) - \hbar\omega/2$. But the energy of the equilibrium fluctuating electromagnetic field is

$$E = \sum_\alpha \hbar\omega_\alpha \left(\frac{1}{2} + \frac{1}{\exp(\hbar\omega_\alpha/kT) - 1}\right),$$

where ω_α are the eigenfrequencies, depending on d.

C) Consider a particle which is at rest in the gravitational field of a Schwarzchild black hole. Its fourvelocity is

$$u^\alpha \equiv \frac{\mathrm{d}x^\alpha}{\mathrm{d}\tau} = \left(\left(1 - \frac{2M}{r}\right)^{-1/2}, 0, 0, 0\right).$$

The proper acceleration of the particle is

$$a^\alpha \equiv \frac{Du^\alpha}{\mathrm{d}\tau} = \frac{\mathrm{d}u^\alpha}{\mathrm{d}\tau} + \Gamma^\alpha_{\beta\gamma} u^\beta u^\gamma = \Gamma^\alpha_{tt} u^t u^t \ (\alpha, \beta, \gamma = t, r, \theta, \varphi)\,.$$

The only nonvanishing component of Γ^{α}_{tt} is

$$\Gamma^{r}_{tt} = \frac{M}{r^2}\left(1 - \frac{2M}{r}\right).$$

Hence $a^{\alpha} = \left(0\,\frac{M}{r^2}, 0, 0\right)$, while

$$|a| = (g_{\alpha\beta}\, a^{\alpha} a^{\beta})^{1/2} = \left(1 - \frac{2M}{r}\right)^{-1/2} \frac{M}{r^2}. \tag{8}$$

A stationary distant observer will measure

$$b^{\alpha} = \frac{Du^{\alpha}}{d\tau}\frac{d\tau}{dt} = a^{\alpha}\left(1 - \frac{2M}{r}\right),$$

$$|b| \equiv (g_{\alpha\beta}\, b^{\alpha} b^{\beta}) = \frac{M}{r^2}. \tag{9}$$

Consequently, the peak of the potential barrier (localized in the vicinity of $r = 3M$) has a nonzero proper acceleration $\cong (3\sqrt{3}\,M)^{-1}$.

D) According to Fabbri (Ref. [5]), who studied the scattering and absorption of electromagnetic waves by a nonrotating black hole, when the frequency w of radiation is smaller than the critical frequency $w_c = (2/27)^{1/2} M^{-1}$, turning points exist for all partial waves, that is for all values of l. When $\omega > \omega_c$, turning points exist only for high $-l$ waves; more precisely, they exist if l is greater than a critical parameter l_c given by the equation $l_c(l_c + 1) = 27\,\omega^2 M^2$.

At high frequencies ($\omega \gg \omega_c$), for $l \ll l_c$, the waves pass above the potential barrier completely unaffected. When l is sligtly greater than l_c the turning points are approximately given by

$$r_{12} = 3M\left\{1 \mp \frac{1}{\sqrt{3}}\left[1 - \frac{27\,\omega^2 M^2}{l(l+1)}\right]^{1/2}\right\}. \tag{10}$$

For $l \leqq l_c$, the zeros of the wave number are given by

$$\bar{r}_{12} = 3M\left\{1 \mp \frac{i}{\sqrt{3}}\left[1 - \frac{l(l+1)}{27\,\omega^2 M^2}\right]^{1/2}\right\}. \tag{11}$$

So, for $\omega > \omega_c$ the transmission coefficient of the barrier is

$$\left.\begin{aligned} T_l &= 0 \quad \text{at} \quad l > l_c \\ T_l &= 1 \quad \text{at} \quad l < l_c \end{aligned}\right\}. \tag{12}$$

For $\omega \ll \omega_c$ real turning points exist for all the partial waves:

$$T_l \cong 4\left[\frac{(l+1)!\,(l-1)!}{(2l)!\,(2l+1)!!}\right]^2 2\,\omega M)^{2l+2}. \tag{13}$$

That is why only the low-frequency waves can successfully escape from the region formed by the Casimir plates with the reflection picture nicely mimiced by (1). Sanchez (Ref. [8]) arrived at a similar conclusion: the reflecting properties of the potential barrier provide that Hawking emission is only significant in the frequency range $0 \leqq \omega < M^{-1}$. Consequently, the potential barrier of a black hole should be approximated by two concentric shells with the first in the vicinity of the horizon ($r = r_1 = 2M + 4\omega^2 M^3/l(l+1) \to 2M$ when $\omega \to 0$) and the second far away from it ($r = r_2 \cong \sqrt{l(l+1)}/\omega$ at $\omega \to 0$). The success of the approximation of the Casimir sphere by two parallel conductor plates (Ref. [2]) permits us to replace each spherical conductor by two plane conductors.

E) Consider an observer resting on surface of a conductor in the gravitational field of a Schwarzchild black hole ($r = r_0$). According to the Principle of Equivalence, he is equivalent to an observer accelerated in Minkowski spacetime with proper acceleration $b^{-1} = \left(1 - \frac{2M}{r_0}\right)^{-1/2} M/r_0^2$. As is well-known (see, for instance [9], and the references cited therein), an observer who is accelerated in flat spacetime with a proper acceleration b^{-1}, should find himself in a thermal bath with temperature $T = b^{-1}/2\pi c k$. The observer who is accelerated with the surface of the wall should find the thermal radiation being in equilibrium with the wall at the same temperature. Thus an observer resting on the surface of a conductor in the gravitational field of a Schwarzchild black hole should discover thermal radiation in equilibrium with a conductor at the temperature

$$T = \frac{M}{r_0^2\left(1 - \frac{2M}{r_0}\right)^{1/2} 2\pi c k}. \tag{14}$$

Hence, *the Interactions of the Radiation with the Surface of the Potential Barrier can be described in Terms of Temperature Corrections to the Casimir Effect.*

(a) The temperature T_1 of a conductor in the vicinity of the horizon is considerably higher than that of a conductor far away from it. So, though each conductor is in equilibrium with radiation, the whole system is in nonequilibrium ($T_1 > T_2$) and a flow of scalar (or any zero-restmass field) establishes itself in the $[r_1, r_2]$ region. The flow is directed from the horizon to spatial infinity.

The observer at rest near the horizon will discover a flow of thermal radiation with a temperature $T_1 = \frac{1}{2\pi}\frac{M}{r_0^2}\left(1 - \frac{2M}{r_0}\right)^{-1/2}$. A distant stationary observer at future infinity J^+ will find that the temperature of radiation in the vicinity of the horizon is $T = M/2\pi r_0^2$. Indeed, the gravitational blue shift of the photon (ratio of observed energy $\hbar\omega_0$ to at J^+ emitted energy $\hbar\omega$) is

$$\omega_0/\omega = (g_{00})^{-1/2} = \left(1 - \frac{2M}{r_0}\right)^{-1/2}.$$

But $\omega/T = \text{const}$ along the light ray (see [10]). That is why $T_1 = T\left(1 - \frac{2M}{r_0}\right)^{-1/2}$.

According to (9), M/r_0^2 is the magnitude of the acceleration (measured by an observer at J^+) of a particle at rest in the gravitational field of a nonrotating black hole. It tends (see [11]) to the so-called "surface gravity" $\varkappa$ when the particle is infinitesimally close to the event horizon. For a Schwarzchild black hole $\varkappa = (4M)^{-1}$ $(c = G = I)$. So, the temperature of the radiation near the horizon is $T_1 = \varkappa/2\pi$ as seen by an observer at J^+. Since the temperature T_2 far from the horizon is negligible, the Poynting vector (see (5), (6)) is

$$P = \int_0^\infty p(\omega)\,\mathrm{d}\omega = \frac{1}{\pi^2}\int_0^\infty \frac{\hbar\omega M(\infty)\,\mathrm{d}\omega}{\exp[\hbar\omega/kT_1] - 1}$$

$$= \frac{\hbar}{\pi^2 c^2}\,\frac{1}{\left(1 - \frac{2M}{r}\right)^2} \cdot \int_0^\infty \frac{\omega_0^3\,\mathrm{d}\omega_0}{\exp[\hbar\omega_0\cdot 2\pi/k\varkappa] - 1}. \qquad (15)$$

The equation obtained exactly coincides with the results of the numerous studies of the Hawking radiation produced on the basis of usual quantum field theory in curved spacetimes. Zero-point oscillations do not contribute directly to the (15) energy flow. But, of course, they influence it through the expression for the energy of the equilibrium fluctuating field

$$E = \sum_\alpha \{\hbar\omega_\alpha + \hbar\omega_\alpha/(\exp[\hbar\omega_\alpha/kT] - 1)\},$$

where the eigenfrequencies ω_α depend on d.

(b) To give a more complete description of the vacuum stress-tensor in the region $[2M, \infty]$ we can use (2)–(4) with d being the "distance" in the accelerated (or Rindler) frame of reference:

$$d = \xi = \left(1 - \frac{2M}{r}\right)^{1/2} r^2/M.$$

But the temperature $T = T(r)$ varies from one point to another. Hence we shall calculate the $\langle T^{\mu\nu}\rangle_{\text{vac.}}$ in the vicinities of r_1 and r_2 first.

The proper acceleration of the r_2 barrier is $b_2^{-1} = \frac{M}{\Delta_2^2}\left(1 - \frac{2M}{\Delta_2}\right)^{-1/2}$, where

$$\Delta_2 = ([l(l+1)]^{-1/2}/\omega - M) \to \infty$$

if $\omega \to 0$. The spherical conductor that is far from the horizon can be represented by two plane conductors with equal temperatures $b_2^{-1}/2\pi$ and accelerations b_2^{-1}. To describe the region $[r_2, \infty]$, the $d \to \infty$ limit of (3) should be involved. Here $(Td) \gg 1$, and $\langle T^{\mu\nu}\rangle_{\text{vac.}}$ is dominated by the Stephan-Boltzmann value all over the space:

$$A = 0, \quad C = 0, \quad B = \frac{\pi^2 T^4}{90} = \frac{M^4}{1440\pi^2\left(1 - \frac{2M}{r}\right)^2}. \qquad (16\,\text{a})$$

To describe the situation near the other side of the r_2 barrier it should be noted that the spherical conductor can be exchanged with a pair of flat plates that rest in the Schwarzchild gravitational field. Applying De Witt's (see [12]) expression for the scalar field vacuum stress-tensor between the Casimir plates, we obtain

$$\langle T^{\mu\nu}\rangle_{\text{vac.}} = \frac{\pi^2}{1440\,d^4}\,\mathrm{diag}(-1, 1, 1, 3) \qquad (16\,\text{b})$$

$$= \frac{\pi^2 M^4}{1440\,r^8\left(1 - \frac{2M}{r}\right)^2}\,\mathrm{diag}(-1, 1, 1, 3).$$

Equations (16) are in good qualitative agreement with Frolov's exact calculations [13] obtained by the usual quantum field methods for Boulware vacuum:

$$\langle T^{\nu}_{\mu}\rangle = \frac{M^2}{1440\,\pi^2\, r^6}\left\{\frac{(2-1.5\,x)^2}{(1-x)^2}(-\,\delta^{\nu}_{\mu} + 4\,\delta^{\nu}_{0}\,\delta^{0}_{\mu}) + 6\,(3\,\delta^{\nu}_{0}\,\delta^{0}_{\mu} + \delta^{\nu}_{1}\,\delta'_{\mu})\right\},$$

$$x \equiv 2\,M/r\,. \tag{17}$$

In the region near the "nonaccelerated" side of the r_2 conductor, (16 b) for negative $\langle T^{\mu\nu}\rangle_{\text{vac.}}$ corresponds to the absence from the Boulware vacuum $|B\rangle$ of blackbody radiation with temperature $T_2 = b_2^{-1}/2\,\pi$. This means that if thermal radiation with T_2 were added, the resulting state would be indistinguishable, near r_2, from the usual Minkowski vacuum. In (16 a) $\langle T^{\mu\nu}\rangle_{\text{vac.}}$ is positive; that corresponds to the presence in the $[r_2, \infty]$ region (near the "accelerated" side of the r_2 conductor) of positive virtual radiation.

The proper acceleration of the r_1 barrier is

$$b_1^{-1} = \frac{M}{(2\,M+\Delta_1)^2}\left(1-\frac{2\,M}{2\,M+\Delta_1}\right)^{-1/2},$$

where $\Delta_1 = 8\,M^3\,\omega^2/l\,(l+1)$.

If we exchange the r_1 conductor with two plates at the distance d apart, we can apply De Witt's equation to describe the situation in the region $[2\,M, r_1]$:

$$\langle T^{\mu\nu}\rangle_{\text{vac.}} = \frac{\pi^2\,M^4}{1440\,r^8\left(1-\dfrac{2\,M}{r}\right)^2}\,\operatorname{diag}(-1,1,1,3)\,. \tag{18a}$$

Or, equivalently, we can use the $(Td) \ll 1$ limit of (2) with $d \cong 4\,M\left(1-\dfrac{2\,M}{r}\right)^{1/2}$:

$$B=0\,,\quad C=0\,,\quad A=\frac{\pi^2}{1440\,d^4} \cong \frac{\pi^2}{1440\,(4\,M)^4\,(1-2\,M/r)^4}\,,$$

$$\langle T_{00}\rangle \cong -\frac{\pi^6}{90\,(8\,\pi\,M)^4\left(1-\dfrac{2\,M}{r}\right)^2}\,. \tag{18b}$$

Again the equations obtained are in good agreement with Frolov's equation (17) for Boulware vacuum. The $|B\rangle$ vacuum in $[2\,M, r_1]$ is depressed below zero by an amount corresponding to the absence from the vacuum of blackbody radiation at a temperature $T = 1/8\,\pi\,M\left(1-\dfrac{2\,M}{r}\right)$.

It is this purely virtual Casimir energy that enables the black hole to contract with nonuniform acceleration.

To describe the situation near the other side of the r_1 conductor, in the direction of acceleration, we can again utilize the $(Td) \gg 1$ limit of (3) with $d \to \infty$. $\langle T_{\mu\nu}\rangle$ is dominated by the Stephan-Boltzmann value all over the conductor:

$$A=C=0\,,\quad B=+\frac{\pi^2\,T^4}{90} = \frac{\pi^2}{90\,(8\,\pi\,M)^4\left(1-\dfrac{2\,M}{r}\right)^2}\,. \tag{19}$$

The expression (16 b) for the "unaccelerated" side of the r_2 barrier can be obtained in a way that clearly reveals its physical significance. If we exchange the r_2 conductor by two plates at distance $\Delta_2\left(1-\dfrac{2\,M}{\Delta_2}\right)^{1/2}$ apart, the $(Td) \ll I$ limit of (2) can be involved to describe the situation in the vicinity of r_2:

$$B=C=0\,,\quad A=\frac{\pi^2}{1440\,d^4} \cong \frac{\pi^2}{1440\left(1-\dfrac{2\,M}{\Delta_2}\right)^2\Delta_2^4}\,,$$

$$\langle T_{00}\rangle = -\frac{\pi^2}{1440\left(1-\dfrac{2\,M}{\Delta_2}\right)^2\Delta_2^4}\,. \tag{20}$$

Of course, the result is too rough to pretend a detailed description but it helps to reveal an important peculiarity of the radiation picture between the r_1 and r_2 conductors: *any observer in the* $[r_1, r_2]$ *region sees two flows of blackbody radiation crossing each other.* The dominating positive flow with $T_1 = (8\,\pi\,M)^{-1}$ comes from the r_1 conductor and the r_2 conductor implies the second one, of negative energy. It comes from the surface of r_2 and corresponds to the absence from the vacuum of blackbody radiation with $T_2 = (2\pi b_2)^{-1}$, according to an observer at J^+. An observer that rests ($r = r_0$) inside the domain $[r_1, r_2]$ sees two flows with

$$T_1 = \frac{1}{8\pi M\left(1 - \frac{2M}{r_0}\right)^{1/2}} \quad \text{and}$$

$$T_2 = \frac{1}{2\pi\xi} = \frac{M}{2\pi r_0^2\left(1 - \frac{2M}{r_0}\right)^{1/2}}.$$

The resulting flow

$$\langle T^{00}\rangle = \frac{\pi^2}{90}[T_1^4 - T_2^4] = \frac{\pi^2(1 - x^8)}{90(8\pi M)^4(1 - x)^2}, \quad x \equiv \frac{2M}{r} \tag{21}$$

is in complete agreement with (7) for $d = 0$. Equation (21) is also in good qualitative agreement with Page's [14] exact formulae obtained for the Hartle-Hawking vacuum:

$$\langle T_\mu^\nu\rangle_\mathrm{H} = \frac{\pi^2}{90(8\pi M)^4} \cdot \left\{\frac{[1 - (4 - 3x)^2 x^6][\delta_\mu^\nu - 4\delta_0^\nu\delta_\mu^0]}{(1 - x)^2} + 24x^6(3\delta_0^\nu\delta_\mu^0 + \delta_1^\nu\delta_\mu^1)\right\}. \tag{22}$$

Thus, all the thermal radiation is "born" within the region $[r_1, r_2]$ between the conductors. Its blackbody spectrum is produced by the interaction of zero-rest-mass fluctuations with the conductor surfaces. The dominating flow is directed from r_1 to r_2 ($T_1 > T_2$). The particles between the conductors are virtual ones. And they would remain virtual if this were the case for real black holes. Yet it is only the scattering aspect of the Schwarzchild gravitational field that was described by our ideal model. The exchange of the black hole potential barrier with two remoted ideal conductors is an approximation merely. The real potential barrier of a black hole forms a "bell" that lasts continuously from zero magnitude at the horizon up to zero magnitude at spatial infinity passing through the maximum at $r = 3M$. The reflecting properties (1), (10)–(13) provide that the barrier behaves as a real conductor and not an ideal one. It conducts well at low frequencies, but as the frequencies increase, its conductivity diminishes. Hence the Hawking radiation is "born" inside the "bell" formed by a potential barrier of a nonrotating black hole in all the region $[2M, \infty]$. The Hawking flow is directed from the $[2M, 3M]$ region to $[3M, \infty]$ tail of the potential barrier. The particles inside the bell are virtual ones. But they can become real after passing through the $[3M, \infty]$ tail, appearing to an observer at future infinity J^+ as "real" ones, created by the accelerated tail of the potential barrier.

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