Hawking-Unruh Thermal Radiance as Relativistic Exponential Scaling of Quantum Noise *

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Abstract

The Hawking-Unruh effect of thermal radiance from a black hole or observed by an accelerated detector is usually viewed as a geometric effect related to the existence of an event horizon. Here we propose a new viewpoint, that the detection of thermal radiance in these systems is a local, kinematic effect arising from the vacuum being subjected to a relativistic exponential scale transformation. This kinematic effect alters the relative weight of quantum versus thermal fluctuations (noise) between the two vacua. This approach can treat conditions which the geometric approach cannot, such as systems which do not even have an event horizon. An example is the case of an observer whose acceleration is nonuniform or only asymptotically uniform. Since this approach is based on concepts and techniques of non-equilibrium statistical mechanics, it is more adept to dynamical problems, such as the dissipation, fluctuation, and entropy aspects of particle creation and phase transitions in black hole collapse and in the early universe.

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1 Introduction

The conventional view of the Hawking-Unruh effect [1, 2] (thermal radiation observed by an accelerated observer [3], from a moving mirror [4], or black hole [4, 5] and for certain observers in an inflationary universe [5, 6, 7]) is based on global properties of spacetime (existence of event horizon), or thermal field theory (periodicity in the propagator) [8, 9]. In 1987 [10] I proposed an alternative approach, viewing it as a local, kinematic effect arising from the vacuum being subjected to an exponential scale transformation. I stressed that the salient features of inflation and black hole collapse are dominated by the infrared behavior of quantum fields, an effect O’Connor and I [11] called ‘dynamical’ finite size effect. Later, with Zhang, I [12] used the analogy with the Kadanoff-Migdal transformation in critical phenomena [13] to describe an inflationary universe, and proposed to view the late stage of inflation and black hole collapse as approaches to the critical regimes [14].

Here, I shall further develop this view by generalizing to relativistic exponential scaling and, using the recently developed theory of noise from quantum fluctuations, show how these effects viewed in one vacuum (Rindler, Hartle-Hawking, or Gibbons-Hawking in the cases of uniformly accelerated detector, black hole and de Sitter universe) can be understood as resulting from the scaling of quantum noise in another (Minkowski). Depicting Hawking-Unruh effect in the light of quantum and thermal fluctuations was first proposed by Sciama [15].

As a pedagogical illustration (mainly for audiences of field theory and statistical mechanics, not general relativity), I present a simple derivation of the de Sitter metric by invoking only the scale transformation concept and special-relativistic effects, and show that Hawking radiation in these cases can be viewed as arising from exponential-redshifting. Though not new, our emphasis on the kinematic aspect over the conventional global geometric or thermal field aspects is intended to clear the conceptual and technical pathways in order to introduce non-equilibrium statistical field theory [16, 17, 18, 20, 22, 23, 24, 25]. In our view this theory based on the use of open system concepts [26] and influence functional methods [27] is more suitable for treating conditions not easily amenable by traditional methods, such as systems which do not have an event horizon. We have started this investigation recently. Analysis of typical cases [28, 29] shows that in such systems radiance indeed is observed, albeit not in an exact Planckian spectrum. The ideas of exponential scaling [10] and results on the infrared behavior of quantum fields [11] can be applied to the analysis of critical phenomena possibly occurring in black holes, early universe and in the semiclassical to quantum gravity transitions.

The stochastic approach we have adopted bears in relation to black hole thermodynamics [30, 31] the same way as stochastic and non-equilibrium statistical mechanics bear to equilibrium thermodynamics. A near-thermal Unruh radiation from an almost-uniformly accelerated detector and a black hole with its modified radiation can be viewed as the linear-response regime, which can be treated by perturbation theory, as distinguished from the fully non-equilibrium conditions of dynamical collapse or arbitrary trajectories, where, we believe, the stochastic theory of particles and fields will prove to be more useful than the conventional methods. (Logically stochastic mechanics connects the foundation, i.e., information and probability theory, to kinetic theory, statistical mechanics and thermodynamics. For an explanation of these conceptual points see [32, 33]). Our approach thus puts more emphasis on statistical mechanics and field theory than on geometry. This has both a technical and
a conceptual rationale. We find the geometric description of spacetime a very elegant, but rather restrictive one, most effective in the large scale, near-equilibrium state of matter and spacetime. I have been of the opinion [34] that Einstein’s theory of general relativity describes only the hydrodynamic or thermodynamic regime of spacetime structure. If we want to probe into the microscopic structure of spacetime and matter, we need to use quantum field theory. (When combined with classical spacetime it defines the semiclassical gravity theory [35] which was instrumental to the discovery of the Hawking-Unruh effect.) If we want to probe into the dynamical and collective properties of spacetime and matter, we need to use non-equilibrium statistical mechanics [36]. Thus we prefer to use the statistical and stochastic quantum field theoretical approach to tackle the new generation of problems like backreaction of particle creation in dynamical collapse, statistical entropy and information loss puzzle in black holes, phase transition in the early universe, and cross-over from semiclassical to quantum gravity. We believe that the usual geometric depiction can be recovered, and is only valid, in the equilibrium and long wavelength (thermo- and hydro-dynamics) limits.

2 Relativistic Exponential Scale Transformation

In this section I will show how exponential scale transformation enters in a central way in the class of spacetimes which admit Hawking-Unruh radiation. To highlight this as a kinematic effect, I’ll use just special relativity to derive the metric of such a spacetime. There is nothing new in the result, which is well-known in general relativity. But this alternative viewpoint and the derivation of an effect usually regarded as belonging to the realm of general relativity could be of some interest for audiences in statistical mechanics and field theory. In this spirit, we will use the ideas of scaling from critical phenomena in conjunction with a quantum theory of noise from stochastic field theory to explore the deeper meaning of these effects and thus prepare the stage for dealing with more complex problems.

2.1 Nonrelativistic scaling

Consider two systems S and Š. S is stationary and Š moves with speed β with respect to S. Let the origins of these two systems coincide at \( t = 0 \), and denote the lengths measured in the two systems as \( r, \tilde{r} \) respectively. They are related by

\[
\tilde{r} = a(t)r \tag{2.1}
\]

where \( a(t) \) is the scale factor relating the two systems. (One can evoke a picture of polka dots on a balloon which is being inflated, such as in Fig. 27.2 of [37], the distance between any neighboring dots increases with time. Indeed one can interchangeably use \( a(t) \) as a measure of time.) Consider the special class of time dependence:

\[
a(t) = e^{Ht} \tag{2.2}
\]

If, say, at each unit time interval \( \Delta t \) the length in \( \tilde{S} \) doubles, i.e., at \( t = 0, \tilde{r} = r; t = 1, \tilde{r} = 2r; t = 2, \tilde{r} = 4r, ... \) Then \( H\Delta t = ln2 \). The exponential transformation is a special
class, because the rate of expansion \( H \equiv \dot{a}/a \) is independent of time. (Compare to a power law \( a(t) = t^p \), \( H = p/t \) depends on \( t \). This seemingly minor point actually makes a great difference in the relation of quantum and thermal fluctuations in these two systems, as we shall see later.) From this we get

\[
H \tilde{r} = Har = \dot{a}r = \beta, \tag{2.3}
\]

which is the relative velocity between the \( S \) and \( \tilde{S} \) systems. It is also the velocity the dots are receding from the origin. Each dot on the expanding balloon sees every other dot moving away from it isotropically, the farther the distance the faster.

From this simple kinematics two related pictures might be evoked:

1) An isotropically expanding universe depicted by the Friedmann-Robertson-Walker model. In particular, if \( a(t) = e^{Ht} \), it is the inflationary universe \([38]\).

2) For the exponential expansion, the Kadanoff-Migdal (KM) transformation.

Recall that in critical phenomena \([13]\) the KM transform is used in conjunction with a block-spin transformation to render a problem defined on a lattice with spacing \( l \) to that of one defined on a larger lattice (after \( n \) iterations) with spacing \([a(\Delta t)]^n l\), and rescaled bond strength. In the above example, \( a(\Delta t) = e^{H\Delta t} = 2\). Near the critical point where many systems manifest scaling behavior, the transformed system preserves the long range characteristics of the original system. This is one way how one could use the ultraviolet behavior of the system to analyze its infrared behavior (via the renormalization group equation, usually constructed from the counterterms introduced for the removal of ultraviolet divergence \([13]\)).

I have used this analogy to explain the scale-invariant properties of inflationary cosmology, and advocated that one can understand the end state of inflation and black hole collapse in terms of the approach to the critical regime, where the infrared behavior of quantum fields in these spacetimes become dominant.

### 2.2 Relativistic Scaling

In the above, the relation \( \tilde{r} = a(t)r \) assumes absolute time, i.e., \( S \) and \( \tilde{S} \) use the same time. Thus the usual KM transformation is non-relativistic scaling. This is of no surprise as there is no concern for relativistic covariance in ordinary critical phenomena studies. Here, to respect special relativity, the two times should be related by a Lorentz factor \( \gamma \), i.e.,

\[
\frac{a(\Delta t)}{a(t)} = \gamma = \frac{1}{\sqrt{1 - \beta^2}} \tag{2.4}
\]

To make precise the terminology, we shall call the following cases defined by Eq. (2.1) as nonrelativistic scale transformation (or scaling, for short), Eq. (2.1) and (2.2) as exponential scale (or KM) transformation, Eq. (2.1) and (2.4) as relativistic scaling, Eq. (2.1) and (2.2) and (2.4) as relativistic exponential scaling, or relativistic KM transformation.

Let us examine the meaning of Conditions (2.3) and (2.4). Note that

i) When \( \beta = 0 \), or \( \gamma = 1 \) i.e., \( H = 0 \), \( a \) a constant, \( \tilde{r} = r \) or \( \tilde{r} = 0 \) (at the origin, \( \tilde{r} = r \),
\[ a() = a(t), t = \text{. The two systems are identical.} \]

ii) When \( \tilde{r} \neq 0 \) or \( H \neq 0, \beta \neq 0 \), then \( a() = \gamma a(t) \). The frequency of radiation \( \tilde{\nu} \) measured in \( \tilde{S} \) is redshifted from \( \nu \) measured in \( S \) by a factor \( \gamma \). Thus \( \gamma \) is the redshift factor. For exponential scaling, \( H = \text{constant} \), we have exponential red-shifting.

iii) When \( \beta \to 1 \) or at \( \tilde{r} = H^{-1}, \gamma \to \infty \). There exists an event horizon of infinite redshift.

Let us now see how the redshift factor can be attributed to a non-flat metric. The easiest way is to recall how gravitational red-shift can be viewed as a special relativistic effect. It is determined by the \( g_{00} \) component of the metric. More precisely,

\[ \tilde{g}_{00} = 1 - \beta^2, \quad \text{or,} \quad \sqrt{\tilde{g}_{00}} = \gamma^{-1}. \]

In the example of a (2D) Schwarzschild metric for a massive (M) object,

\[ ds^2 = \left(1 - \frac{2M}{r}\right)dt^2 - \frac{dr^2}{\left(1 - \frac{2M}{r}\right)}, \]

we can see the following analogous cases:

i) At \( r = \infty, g_{00} = 1, \gamma = 1 \), the metric is asymptotically flat.

ii) At \( r > 2M, 1 > g_{00} > 0, \gamma > 1 \).

iii) At \( r = 2M, g_{00} = 0, \gamma \to \infty \). This infinite-redshift surface is, of course, what defines the black hole horizon.

Hence by analogy, we see that for the coordinates \((\tilde{r}, \tilde{\xi})\) in the \( \tilde{S} \) frame, \( \tilde{g}_{00} = \gamma^{-2} = 1 - (H\tilde{r})^2 \). The metric embodying an exponential scale transformation is thus of the form:

\[ ds^2 = [1 - (H\tilde{r})^2]d\tilde{t}^2 - \frac{d\tilde{r}^2}{[1 - (H\tilde{r})^2]} \]

This is the form of a de Sitter metric in the so-called static coordinate. Note that an event horizon exists at \( \tilde{r} = H^{-1} \) for observers at \( \tilde{r} = 0 \). The exponential red-shifting factor from relativistic scaling is what we prefer to focus on (rather than the geometric structure) in the statistical field theory approach to thermal effects and their generalizations.

### 3 Thermal Particle Creation from Exponential Red-shifting

Particle creation from a black hole (Hawking effect) or observed by a uniformly accelerated observer (Unruh effect) have been treated in many ways by methods in quantum field theory in curved spacetimes [35]. Of the many existing derivations and interpretations, I want to highlight what I see as the central aspect of this problem, i.e., the vacuum state where thermal radiance is observed (e.g., the Rindler vacuum of a constantly accelerated detector in Unruh effect) is related to the inertial vacuum state (Minkowski vacuum) by an exponential scale transformation. This transformation is of the same form and nature as that between the \( S, \tilde{S} \) systems in the simple example given above. In fact, the radiation observed by an observer in the \( \tilde{S} \) system is also of a thermal nature. It was first deduced by Gibbons and
Hawking [3] by analogy with the Hawking radiation in a Schwarzschild metric. In all three cases, i.e., the Rindler, the Schwarzschild and the de Sitter metrics, it is the exponential redshifting which is responsible for the thermal nature of radiation. We make the fine distinction that it is the exponential redshifting and not the existence of an event horizon which is the necessary condition. The above situations encompass both the cases where there is an event horizon but no (thermal) particle creation (like an extreme Reisner-Nordström metric), and cases where the horizon is not globally defined (like the case of finite-time acceleration). The former case was brought up in a discussion between me and Unruh during his visit to Maryland in 1988 but never pursued. The latter case is treated by Raval, Koks and myself [28] recently, using the statistical field theory methods.

Consider the case of a uniformly accelerated detector. If we write the Minkowski metric in the null coordinates \((U, V)\)

\[
ds^2 = dt^2 - dx^2 = dUdV \tag{3.1}
\]

where

\[
U = t - x, \quad V = t + x, \tag{3.2}
\]

and perform a conformal transformation to the Rindler coordinates \((\xi, \eta)\),

\[
ds^2 = e^{2a\xi}(d\eta^2 - d\xi^2) \tag{3.3}
\]

with the associated null coordinates \((u, v)\)

\[
u = \eta - \xi, \quad v = \eta + \xi, \tag{3.4}
\]

the two sets of null coordinates are then related by

\[
U = - \frac{1}{a}e^{-au}, \quad V = \frac{1}{a}e^{av}. \tag{3.5}
\]

Particle detector moving at constant acceleration \((\xi = \text{const})\) has a trajectory \(x^2 - t^2 = \alpha^2\), where \(\alpha^{-1} = ae^{-ak}\) is the proper acceleration. The detector’s proper time is \(\tau = e^{a\xi}\eta\). (See discussion in, e.g., Sec. 4.5 of [35])

The Schwarzschild metric (2.6) depicting an eternal (2D) black hole can be written in terms of the Regge-Wheeler coordinates

\[
r^* = r + 2M\ln|\left(\frac{r}{2M}\right)| - 1| \tag{3.6}
\]

as

\[
ds^2 = (1 - \frac{2M}{r})(dt^2 - dr^*^2) = (1 - \frac{2M}{r})dudv, \tag{3.7}
\]

where \((u, v) = t - r^*, v = t + r^*\) are the null Schwarzschild coordinates. In analogy with the uniformly accelerated observer case, one can introduce a set of Kruskal coordinates, \((\bar{t}, \bar{r}^*)\), and write the metric as (see, e.g., Sec. 3.1 of [35])

\[
ds^2 = \frac{2M}{r}e^{-\frac{2M}{r}}(dt^2 - dr^*^2) = \frac{2M}{r}e^{-\frac{2M}{r}}(dUdV) \tag{3.8}
\]
The null Schwarzschild and Kruskal coordinates are related by

\[ U = -4Me^{-\frac{4}{v}} , \quad V = 4Me^{\frac{4}{v}} \]  \hspace{1cm} (3.9)

Note the pairwise correspondence between the Minkowski / Kruskal \((U, V)\) and the Rindler / Schwarzschild \((u, v)\) coordinates in the accelerated detector and the black hole cases. Note again the exponential relation between these two sets of coordinates.

Let me briefly describe the relation between Killing vectors, normal modes, vacuum states, Bogolubov transformation and particle creation in a curved spacetime. (The reader is referred to e.g., [35] for details). The existence of a Killing vector in a spacetime (e.g., \(\partial_t\), or equivalently, \(\partial_U, \partial_V\) in Minkowski space) allows for a normal mode decomposition. The amplitudes of the normal modes when second quantized define the creation and annihilation operators \((A, A^+)\), and the number operator \(n = A^+A\) which make up the Fock space with respect to this decomposition. The vacuum \(|0, t\rangle\) being the no particle state is defined by taking modes to be positive frequency with respect to the Killing vector \(\partial_t\). The amplitude functions or the annihilation and creation operators \((a_j, a_j^+)\) of another set of modes decomposed with respect to another Killing vector (e.g., \(\partial_u\)) is related to the original one by a Bogolubov transformation:

\[ a_j = \Sigma_i (\alpha_{ij} A_i + \beta_{ij}^* A_i^+) \]  \hspace{1cm} (3.10)

where \(\alpha, \beta\) as the Bogolubov coefficients of the \(kth\) mode.

In the case of a collapsing mass, an incoming wave from past infinity \(I^-\) in the form \(e^{-i\omega v}\) (the Killing vectors \(\partial_u, \partial_v\) define the Schwarzschild or the Boulware vacuum) falling towards the mass would be subjected to a blue-shift. Having passed through the collapsing mass, the outgoing wave while climbing out of the severe and increasing gravitational potential is subjected to an exponential redshifting which far exceeds the blue-shift (there is also a small spin-dependent contribution from the effective potential), i.e.,

\[ e^{-i\omega v} \rightarrow e^{+i\omega U} = e^{+i\omega(-4Me^{-\frac{4}{v}})} \]  \hspace{1cm} (3.11)

It is seen that the exponential arises from the defining relation of \((u, v)\) and \((U, V)\), (the Killing vectors \(\partial_U, \partial_V\) are used in the case of an eternal black hole – black hole in equilibrium with its Hawking radiation – to define the Hartle-Hawking vacuum). The so-called Unruh vacuum (defined with an ‘in’ state with respect to \(\partial_v\) and ‘out’ state with respect to \(\partial_U\)) is most suitable for the description of the actual black hole collapse and radiation emittance situation.

For the above set-up, the coefficient \(\beta\) connecting a positive frequency incoming component \(e^{-i\omega v}\) and a negative frequency outgoing component \(e^{+i\omega(-4Me^{-u/4M})}\) has the special form:

\[ |\beta_k/\alpha_k|^2 = e^{-8\pi M \omega} \quad \text{or} \quad e^{-2\pi \omega/a}, \]  \hspace{1cm} (3.12)

where the second term is for the accelerated observer. This leads to a Planckian spectrum

\[ < n_k > = (e^{8\pi M \omega} - 1)^{-1} \quad \text{or} \quad (e^{2\pi \omega/a} - 1)^{-1}, \]  \hspace{1cm} (3.13)

which gives respectively the Hawking and Unruh temperatures \(T_H, T_U\)

\[ k_B T_H = \frac{1}{8\pi M} , \quad k_B T_U = \frac{a}{2\pi}, \]  \hspace{1cm} (3.14)
where \( k_B \) is the Boltzmann constant. Note that this form is readily identifiable from the exponential form of the relations (3.5) (3.9) between \((u, v)\) and \((U, V)\) defining the in and out states. The out wave being exponentially redshifted is responsible for the thermal nature of the Hawking-Unruh radiation.

Let us now return to the observers in \( S \) and \( \tilde{S} \). The two vacuum states \(|0>_{s}, |0>_{\tilde{s}}\) defined with respect to the two Killing vectors \( \partial_t, \partial \) bear the same relation as the Schwarzschild versus the Kruskal vacuum. By analogy, an observer in \( \tilde{S} \) will therefore see a thermal radiation with temperature given by

\[
k_B T_{ds} = \frac{H}{2\pi}.
\]

(This was first derived by Gibbons and Hawking [5].) We use the de Sitter space for illustration because it manifests the kinematic effect most directly. But for all three cases, we can say that thermal radiance detected by one observer (in \( \tilde{S} \) ) arises from the relativistic exponential scaling of vacuum fluctuations of the other (in \( S \)).

To complete our thesis, it remains to show how vacuum fluctuations can be understood as quantum noise. For this we need to introduce some basic notions and techniques of statistical field theory. We will try to illustrate the basic ideas with minimal technical detail.

4 Quantum Noise under Exponential Scaling Manifests as Thermal Radiation

Noise from vacuum fluctuations [24] can best be defined in terms of the stochastic theory of quantum fields using the quantum open system concept and the influence functional formalism. A general review of this method is given in [22]. Details of the following summary are contained in [18].

We consider an Unruh-DeWitt detector undergoing constant acceleration \( a \) with trajectory (in 2D)

\[
x(\tau) = \frac{1}{a} \cosh a \tau, \quad s(\tau) = \frac{1}{a} \sinh a \tau
\]

(4.1)

where \((x, s)\) are its internal coordinates and \( \tau \) its proper time. For a 2D scalar quantum field \( \Phi(x_0, \eta) \) in flat space, one can decompose it in normal modes, each describable by a Lagrangian

\[
L(s) = \frac{1}{2} \sum_{\sigma} \sum_{k} [(q_{k}^{\sigma})^2 - k^2 (q_{k}^{\sigma})^2]
\]

(4.2)

where \( q_{k}^{\sigma} \) is the amplitude functions for the kth mode (here \( \sigma = \pm \) denotes the sin and cos standing wave components). The detector-field interaction is described by an interaction Lagrangian density

\[
L_{int}(x) = -\epsilon r \Phi(x) \delta(x(\tau)).
\]

(4.3)

The influence of the quantum field on the detector is expressed in terms of an influence kernel which has the form

\[
\zeta(s(\tau), s(\tau')) = \nu(s, s') + i \mu(s, s')
\]

(4.4)
where $\nu, \mu$ are the noise and dissipation kernels. For an uniformly accelerated observer it is found by Anglin, Hu and Matacz \[17, 18\] to be,

$$
\zeta(\tau, \tau') = \int_{0}^{\infty} dk I(k) \left[ \coth(\pi k/a) \cos k(\tau - \tau') - i \sin k(\tau - \tau') \right],
$$

(4.5)

where $I(k)$ is the spectral density function. By comparison with the Brownian motion model \[27\], it is seen to have the same form as a particle in a thermal bath with temperature given by (3.14), the expected Unruh temperature. Similarly, for a 2D black hole, we see that a detector at $\bar{r}^* = \text{constant}$ is influenced by the quantum field in the same way (same form of the influence functional) as an inertial detector in flat 2D spacetime; while a detector at $r^* = \text{constant}$ has the same influence functional as an uniformly-accelerated detector in flat 2D spacetime. Using this analogy, it is easy to see that the influence kernel of a scalar field on a detector at $r^* = \text{constant}$ (which defines the Hartle-Hawking vacuum) has the form

$$
\zeta(t, t') = \int_{0}^{\infty} dk I(k) \left[ \coth(4\pi M k) \cos k(t - t') - i \sin k(t - t') \right],
$$

(4.6)

By comparison with a Brownian oscillator in a thermal bath, it is easily seen that the detector will see a thermal radiation with the Hawking temperature (3.14).

This derivation of the Hawking-Unruh radiation relies mainly on the statistical field rather than geometric ideas. In this approach, it is seen that the quantum noise in one set of vacuua (the Minkowski, Kruskal or the S observer) becomes thermal radiation in another set of vacuua (the Rindler, Schwarzschild or the $\tilde{S}$ observer). The latter is related to the former by a relativistic exponential scale transformation, the exponential red-shifting being responsible for the exact thermal but coherent characteristics of the radiance.

One may think that the kinematic and statistical approach we have presented here is just another description equivalent to the traditional geometric approach. This is true for exactly thermal (equilibrium) conditions, corresponding to the special case of uniform acceleration. But we think the statistical field theory approach has an advantage over the geometric approach in treating cases which deviate from these conditions, e.g., for detectors which undergo acceleration for only a finite interval of time, or approach uniform acceleration asymptotically, as well as dynamical collapsing mass, and cosmological models with near-exponential expansion. In these cases there does not exist an event horizon, so the traditional arguments depending on such a condition (such as the periodicity in the thermal propagator) would not be available. However, one can still use concepts such as near-exponential scale transformation and perturbation techniques in field theory to treat these cases. In two recent papers, Raval, Koks, Matacz and I \[28, 29\] have shown that in such systems radiance indeed is observed, albeit not in a precise Planckian spectrum. The deviation therefrom is determined by a parameter which measures the departure from uniform acceleration or from exact exponential expansion. These results are expected to be useful for investigating the non-equilibrium black hole thermodynamics and the linear-response regime of quantum backreaction problems.

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