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The Black Hole Event Horizon as a Limited Two-Way Membrane

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The Black Hole Event Horizon as a Limited Two-Way Membrane

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I. Introduction

Black holes radiate [1]. This Hawking radiation is generated from sourceless virtual vacuum field fluctuations of the vacuum state of spacetime [2-6,21]. The radiated virtual particles transmute into real particles [2-6].

That black holes radiate leads to the contemplation that perhaps the classical wisdom that nothing can escape from a black hole is somewhat superficial [7]. One might reasonably speculate that other kinds of virtual vacuum field fluctuations exist which can escape from a black hole as well.

Virtual vacuum field fluctuations do also exist as exchange particles, which constitute the mechanism of quantum field interactions in spacetime and are governed by the rules of quantum field theory [13-17,20,22]. Since the local geometry of spacetime seamlessly continues as one crosses the event horizon of a black hole [2,3,5,9,28], quantum field interactions continue in the crossing [2,3,5,8-9]. Messages and matter (and so virtual particle exchanges) can be sent from sources outside a black hole to sinks inside [3]. Such quantum field interactions persist among sources and sinks in a black hole's interior [2,3,5,9].

Virtual vacuum exchange fluctuations differ from Hawking radiation. Being generated by a source is one characteristic distinguishing these virtual exchange particles from the sourceless virtual particles of Hawking radiation. Further unlike Hawking radiation, exchange particles remain virtual and are not subject to direct measurement.[1]

It is natural then to enquire as to whether a virtual exchange particle can propagate from the interior or event horizon of a black hole to its exterior. This paper proffers a theoretical basis for such a limited phenomenon.

II. Analytic Components

It is necessary to briefly introduce and discuss certain aspects of the principal objects used in this paper's analytic.

a) Schwarzschild black hole

For simplicity, we consider an uncharged, spherically symmetric, non-rotating black hole - a Schwarzschild black hole - as the background spacetime. The black hole's event horizon is a null surface of radially outgoing real photon world paths at radius $R_s = 2GM/c^2$ [3,5,9,12], with radius of curvature at the event horizon

[1] Ref. [14], Sec. 8.4.2-8.4.4.

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given by $\kappa^{-1} \equiv \rho_{eh} = 2R_s$ [9]. The event horizon is considered to be a one-way membrane delineating the boundary from which nothing can escape.[2] No event on or within the event horizon can send a "signal" out to an external region.[3] This is an effect of the null cones of spacetime having null geodesics tangent to the event horizon, thereby permitting matter and signals to pass inwards only [2,3,5,21]. At the event horizon the escape velocity from the black hole becomes the speed of light *in vacuo* [6,12,21,26-27]. The following terminology will be used:[4]

Σ_0 designates the boundary (event horizon) of the black hole;
β^-designates the interior region of the black hole;
β^+designates the exterior region of the black hole;
x_{eh} represents a spacetime event located on Σ_0;
y_+ represents a spacetime event located in β^+;
x_- represents a spacetime event located in β^-;
h designates the r-coordinate distance from the black hole's singularity.[5]

b) Massless real scalar quantum field

We also consider for simplicity a massless real scalar quantum operator field (massless spin-0 field) with mode expansion given by [4,14,16,17]

$$\varphi(\mathbf{x}, t) = \int \frac{d^3\mathbf{k}}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} [e^{-ikx} \hat{a}^-_{\mathbf{k}} + e^{ikx} \hat{a}^+_{\mathbf{k}}], \tag{1}$$

where $\hat{a}^-_{\mathbf{k}}$ and $\hat{a}^+_{\mathbf{k}}$ are the time-independent destruction and creation operators, respectively, and $kx = \omega_{\mathbf{k}} t - \mathbf{k} \cdot \mathbf{x}$. Each $\varphi_{\mathbf{k}}(x, t) = \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} [e^{-ikx} \hat{a}^-_{\mathbf{k}} + e^{ikx} \hat{a}^+_{\mathbf{k}}]$ is a mode field operator. Each $\varphi_{\mathbf{k}}(t) = \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} e^{\pm i\omega_{\mathbf{k}} t}$ is a mode function.

Eq. (1)'s integration over wave vectors $\mathbf{k}$ is equivalent to integration over momentum states $\mathbf{p}$ or wavelengths $\lambda_{\mathbf{k}}$, using the de Broglie relations [13,16,18][6]

$$\mathbf{p} = \hbar\mathbf{k}; \quad \mathbf{p} = h/\lambda_{\mathbf{k}}; \quad \mathbf{k} = 2\pi/\lambda_{\mathbf{k}}; \quad E_{\mathbf{k}} = \hbar\omega_{\mathbf{k}}. \tag{2}$$

The eigenmode $\varphi_{\mathbf{k}}(\mathbf{x}, t)$ has momentum $\mathbf{p} = \hbar\mathbf{k}$ [15,16]. Since each $\mathbf{p}$ gives the specific momentum of a single mode $\varphi_{\mathbf{k}}(\mathbf{x}, t)$ of the quantum field $\varphi(\mathbf{x}, t)$ [16],[7] it thus specifies the mode's wave vector $\mathbf{k}$ with associated unique wavelength $\lambda_{\mathbf{k}}$ and angular frequency $\omega_{\mathbf{k}}$ [14].

Quantum field theory requires the $\omega_{\mathbf{k}}^{-1/2}$ factor of Eq. (1) in order to maintain relativistic covariance of the field description [16,23].[8] The amplitude for occurrence of a quantum fluctuation of a mode $\varphi_{\mathbf{k}}(\mathbf{x}, t)$ is given by [4][9]

$$\delta\phi_{\mathbf{k}} \sim \omega_{\mathbf{k}}^{-1/2}. \tag{3}$$

Thus $\delta\phi_{\mathbf{k}} \to 0$ as $\omega_{\mathbf{k}} \to \infty$. Since $\omega_{\mathbf{k}} \to \infty$ as $\mathbf{k} \to \infty$, it follows from Eq. (2) that $\delta\phi_{\mathbf{k}} \to 0$ as $\lambda_{\mathbf{k}} \to 0$. As wavelength decreases the amplitude for occurrence of a quantum fluctuation of the associated $\lambda_{\mathbf{k}}$-field mode decreases.

c) Vacuum state

The vacuum state $|0\rangle$ is defined as that quantum state with the lowest possible energy [4,8,10], and is taken to satisfy the equation [4,13,14,16]

$$\hat{a}^-_{\mathbf{k}} |0\rangle = 0. \tag{4}$$

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[2] Ref. [6], p. 15.
[3] Ref. [6], pp. 27, 29, 39.
[4] See Ref. [2], Sec. 34.2-4 & Ref. [9], Ch. 12, for detailed technical analysis of these terms.
[5] See Ref. [3] for definition of the term "r-coordinate".
[6] Ref. [16], Eq. 2.47 & Ref. [18], Eqs. 5.4-5.5.
[7] Ref. [16], Sec. 2.4.
[8] Ref. [23], pp. 54-6, 86; Ref. [16], pp. 22-3. RQM requires the same factor to maintain unity of probability - Ref [14], pp. 46-7, 63.
[9] Ref. [4], Sec. 1.4 & 4.4-4.5.

Various issues arise when attempting to define a unique vacuum state in a general curved spacetime [4,8-11,21,29-30]. For instance, in a general spacetime it is considered that no analogue of a positive frequency subspace exists and the notion of a unique time-translation operator or time parameter is not well-defined [11,21,29].[10] This lack of a positive frequency subspace leads to ambiguities regarding frequency mode decomposition of a quantum operator field such as $\varphi(\mathbf{x},t)$ [1,4,8,11,29,30].

Despite these issues, a necessary condition to permit interpretation of $|0\rangle \equiv |0_M\rangle$ as the unique vacuum state of minimum excitation energy in a region of spacetime can be framed [4,11].[11] Framing this condition begins with a quantization procedure for the quantum field $\varphi(\mathbf{x},t)$.[12] We first put the field expansion

$$\varphi(\mathbf{x},t) = \int \frac{d^3\mathbf{k}}{(2\pi)^{3/2}} \frac{1}{\sqrt{2}} [\nu^*_{\mathbf{k}}(t)\, e^{i\mathbf{k}\cdot\mathbf{x}} \hat{a}^-_{\mathbf{k}} + \nu_{\mathbf{k}}(t)\, e^{-i\mathbf{k}\cdot\mathbf{x}} \hat{a}^+_{\mathbf{k}}], \tag{5}$$

with the $\nu_{\mathbf{k}}(t)$ mode functions to be determined. Postulating the quantum commutation relations $\left[\hat{a}^-_{\mathbf{k}}, \hat{a}^+_{\mathbf{k}'}\right] = \delta(\mathbf{k}-\mathbf{k}')$ and using Eq. (4) along with the normalization conditions $\dot{\nu}_{\mathbf{k}}(t)\, \nu^*_{\mathbf{k}}(t) - \nu_{\mathbf{k}}(t)\, \dot{\nu}^*_{\mathbf{k}}(t) = 2i$ on the mode functions $\nu_{\mathbf{k}}(t)$ [4], it is found that $|0\rangle$ can be interpreted as the unique vacuum state only if the $\nu_{\mathbf{k}}(t)$ of Eq. (5) take the form

$$\nu_{\mathbf{k}}(t) = \frac{1}{\sqrt{\omega_{\mathbf{k}}}} e^{i\omega_{\mathbf{k}} t} \propto \varphi_{\mathbf{k}}(t)\,. \tag{6}$$

This condition on the mode functions for use of $|0\rangle$ thus also generates Eq. (1) for the form of the quantum field $\varphi(\mathbf{x},t)$. The generated subspace with well-defined positive frequency modes and quantum operator field permits formulation of a propagation mechanism for a well-defined set of virtual exchange modes.

d) Virtual exchange particles

What we term virtual exchange particles are actually quantum superposed virtual propagating waves [25], or put another way - propagating virtual vacuum field disturbances [14,15,17,25]. This becomes evident when considering the quantum field's action on the vacuum state: $\varphi(\mathbf{x},t)\,|0\rangle$, which excites the vacuum creating a virtual fluctuation or disturbance [14,15,17]. This virtual vacuum disturbance is comprised of an outward propagating continuous superposition of an infinite number of $\mathbf{k}$-eigenstates [14-17], each eigenstate associated with a unique wavelength $\lambda_{\mathbf{k}}$ via the de Broglie relations of Eq. (2) [14-17,22,25].[13]

This virtual exchange particle propagates to a separate spacetime event where it is absorbed [14-17,22]. Of critical import is that though the virtual exchange particle is off its own mass shell [14,15],[14] it must remain on-shell of the dispersion relation related to the mass that it carries from source to sink [14,15,20-22].[15]

To clarify, a source at a spacetime event $\mathbf{X}$ emits a virtual exchange particle Q of mass m_Q which propagates to event $\mathbf{Y}$ where it is absorbed by a sink. The source particle loses mass m_Q, and the sink particle gains this equivalent mass

[10] Ref. [9], Sec. 14.2; Ref. [21], Sec. 24.3 & 30.4.
[11] Where $|0_M\rangle$ designates the vacuum state for a fiat (Minkowski) spacetime region.
[12] See Ref. [4], Sec. 4.3-4.4 for full rendition of this quantization procedure.
[13] Ref. [15], Ch 1.4; Ref. [17], pp. 47-8.
[14] See Ref. [21], Sec. 26.7 for a discussion.
[15] Ref. [22], Sec. 17.4; Ref. [14], Box 8-1. Closed-loop (self-energy) Feynman diagrams and their processes are not being considered herein (which would permit such off-shell propagation), as these processes are not applicable to this paper. See Ref. [14], Sec. 8.4.2, 8.4.5-8.4.6 regarding the virtual photon and closed-loop analysis; Also Ref. [16], Sec. 6.2.

[22]. Thus Q must maintain mass m_Q throughout the propagation, and energy conservation at the vertices of the associated Feynman diagram is given by the dispersion relation [14,22]

$$E_Q = (\mathbf{p}_Q^2 + m_Q^2)^{1/2}. \tag{7}$$

Since Q is represented by a massless scalar field, it must remain off its own mass shell: $E_\mathbf{p} \neq |\mathbf{p}|$, or $p_\mu p^\mu \neq 0$.[16] At the same time since energy is conserved at the vertices of the Feynman diagram, Eq. (7) must be maintained. In short, in Feynman diagrams without closed loops the virtual exchange particle four-momenta are pinned down at each vertex [13,14]; thus $p_\mu p^\mu = m_Q^2$, and the virtual exchange particle Q must abide by Eq. (7)'s mass shell relation.

The virtual exchange particle comprises a superposition over all values $\mathbf{k} = 2\pi/\lambda_\mathbf{k}$ via Eq. (1) [15]. As the wave vector $\mathbf{k}$ (and hence momentum $\mathbf{p} = \hbar\mathbf{k}$) increases in the superposition of field modes, each mode's angular frequency simultaneously increases in order to maintain Eq. (7)'s constraint. This constraint can thus also be written as $\omega_\mathbf{k} = (\mathbf{k}^2 + m_Q^2)^{1/2}$.

e) Feynman virtual exchange particle propagator

The Feynman propagator governing the propagation of virtual exchange vacuum field fluctuations arises in the course of deriving quantum field interaction theory [14,15]. The amplitude $\langle 0| \varphi(\mathbf{y},t)\varphi(\mathbf{x},t) |0\rangle$ concerns only half of the Feynman propagator [16]. The total Feynman propagator is [14,17,22,31-32][17]

$$i\triangle_F(x \mid y) = \Theta\left(y^0 - x^0\right) \langle 0| \varphi(\mathbf{y},t)\varphi(\mathbf{x},t) |0\rangle \tag{8}$$

$$+\Theta\left(x^0 - y^0\right) \langle 0| \varphi(\mathbf{x},t)\varphi(\mathbf{y},t) |0\rangle ,$$

which can be written as

$$i\triangle_F(x \mid y) = \Theta\left(y^0 - x^0\right) i\triangle_F^+(y/x) + \Theta\left(x^0 - y^0\right) i\triangle_F^-(x/y). \tag{9}$$

The Feynman propagator $i\triangle_F(x \mid y)$ corresponds to a single virtual vacuum field excitation propagation [13-17]. It consists of the summations of amplitudes representing the possibility that the propagating virtual exchange disturbance can be either particle $i\triangle_F^+(y/x)$ or antiparticle $i\triangle_F^-(x/y)$ in character [17,22]. The virtual exchange particle propagation only is given by

$$i\triangle_F^+(y/x) = \langle 0| \varphi(\mathbf{y},t)\varphi(\mathbf{x},t) |0\rangle . \tag{10}$$

III. Core Propositions

The foregoing objects and following set of propositions form the core of the technology from which the paper's hypotheses are deduced.

a) Proposition 1

The weak equivalence principle applies to virtual exchange particles.

A standard statement of the weak form of the principle of equivalence of general relativity is that the gravitational field is coupled to all matter and energy [5].

[16] Otherwise, the Feynman propagator threatens to blow up [13-17].

[17] See Ref. [17], Eq. 2.81: with the more appropriate attribution being the *Stueckelberg- Feynman* propagator; Ref. [22], Sec. 6.3 & Eq. 17.29.

This notion needs generalized to include virtual fields. Consider the action for the scalar field in a curved spacetime [4,15]:

$$S=\int d^4x\sqrt{-g}\frac{1}{2}\left(g^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi-V(\varphi)\right). \tag{11}$$

This action depends explicitly on the metric and couples the scalar field to the spacetime metric field [4].

The scalar operator field's operation on the vacuum thereby couples the emitted virtual exchange particle to the gravitational field. This coupling necessarily follows from the principle of general covariance.[18] The coupling of virtual particles to the gravitational field is also evident when considering the mechanism behind the Hawking effect [1,5,9,15,21].

b) *Proposition 2*

For a set of wavelengths $\{\lambda_{\mathbf{k}}:\lambda_{\mathbf{k}}=2\pi/\mathbf{k}\}$ of a quantum field's $\varphi(\mathbf{x},t)$ mode expansion with $\lambda_{\mathbf{r}}\in\{\lambda_{\mathbf{k}}\}$ iff $\lambda_{\mathbf{r}}\ll\rho_{eh}$, a unique vacuum state $|0\rangle\equiv|0_M\rangle$ exists in the region of spacetime Ω comprising the black hole's event horizon Σ_0 and a region exterior to Σ_0.

A basic tenet of general relativity is that on sufficiently small scales curved spacetime can be treated as locally flat [2,4,9,11,28], and in such a region $|0\rangle$ can be associated with the vacuum state $|0_M\rangle$ [4,8,29,30]. As previously noted, in a general curved spacetime defining $|0\rangle$ poses problems [4,8,10,11]. Nonetheless, this general relativistic tenet can be imported into quantum field theory by imposing certain constraints on the local spacetime curvature and set of field modes being considered [4,11,29-30]. In such an arena a positive frequency subspace exists and decomposition of field modes into positive and negative frequency modes is viable [4,8,21].

To do this we first define the region $\mathbf{\Omega}$ of spacetime being considered. Because of spherical symmetry we can consider an infinitesimal symmetric region $\mathbf{\Omega}_1$ emanating off of Σ_0 and into β^+, with $\Sigma_0\subset\mathbf{\Omega}_1$. Designate Σ_1 as the outer boundary of $\mathbf{\Omega}_1$, with $\Sigma_1\subset\mathbf{\Omega}_1$.Thus $\mathbf{\Omega}_1$ is bounded by the hypersurfaces $[\Sigma_0,\Sigma_1]$.

This process can be iterated n times until an outer boundary Σ_n is reached for which $y_+\in\Sigma_n$, and thus $y_+\in\mathbf{\Omega}_n$. The region $\mathbf{\Omega}$ is defined as the covering union of these subregions:

$$\mathbf{\Omega}=\bigcup_{i=1}^{n}\mathbf{\Omega}_i, \tag{12}$$

with this cover bounded by the hypersurfaces $[\Sigma_0,\Sigma_n]$.

The background spacetime being asymptotically flat, it follows that $\rho\to\infty$ as $h\to\infty$. Thus $\forall k<n$ we have $\rho_{\Sigma_n}>\rho_{\Sigma_k}$ and $\rho_{\mathbf{\Omega}_n}>\rho_{\mathbf{\Omega}_k}$. More specifically $\rho_{y_+}>\rho_{eh}$, with $y_+\in\mathbf{\Omega}_k$ & $y_+\notin\Sigma_0$, $\forall k=1\to n$.

Thus $\forall\lambda_{\mathbf{r}}\in\{\lambda_{\mathbf{k}}\}$ it follows that

$$\lambda_{\mathbf{r}}\ll\rho_{\mathbf{\Omega}}, \tag{13}$$

where $\rho_{\mathbf{\Omega}}$ designates the radius of curvature at any spacetime point $x_0\in\mathbf{\Omega}$.

Next, the metric is expanded about $x_0\in\mathbf{\Omega}$ in Riemann coordinates [11]

$$g_{\mu\nu}(x)=\eta_{\mu\nu}+\frac{1}{3}R_{\mu\nu\alpha\beta}(x_0)(x-x_0)^\alpha(x-x_0)^\beta+O[(x-x_0)^3]. \tag{14}$$

[18] Ref. [4], Sec. 5.2.

If the $\mathbf{k}^2$ of Eq. (2) are sufficiently large in comparison with the components of the Riemann tensor $R_{\mu\nu\alpha\beta}$ in a local spacetime region, then a vacuum state $|0_M\rangle$ on flat spacetime can be well-defined in this region for these $\mathbf{k}$-modes [11]. A particle detector will respond to Fock states as it would in Minkowski spacetime, and we thus may put $|0\rangle \equiv |0_M\rangle$ for these modes [11].

This condition is identical to the requirement that the corresponding set of field mode wavelengths $\{\lambda_\mathbf{k}\}$ be sufficiently small when compared with $\rho_\mathbf{\Omega}$ [4]; in particular, we require $\forall\lambda_\mathbf{r} \in \{\lambda_\mathbf{k}\}$ that they possess the property $\lambda_\mathbf{r} \ll \rho_{eh}$ [4,11], which in turn mandates Eq. (13).

Such a construction permits the $\mathbf{k}$-mode functions which are in one-to-one correspondence with the elements of $\{\lambda_\mathbf{k}\}$ to be defined in $\mathbf{\Omega}$ as [4]

$$\varphi_\mathbf{k}(t) \approx \frac{1}{\sqrt{2\omega_\mathbf{k}}} e^{\pm i\omega t}, \tag{15}$$

leading to the definition of $|0\rangle \equiv |0_M\rangle$ as the "in" vacuum state for $\{\lambda_\mathbf{k}\}$ in $\mathbf{\Omega}$ [4], as well as permitting Eq. (1) to represent the form of the quantum field for such modes [4,11].[19] Since $y_+ \in \mathbf{\Omega}_k$ & $y_+ \notin \Sigma_0 \Longrightarrow \rho_{y_+} > \rho_{eh}$, it follows that $\langle 0| \equiv \langle 0_M|$ is a well-defined "out" vacuum state at y_+ for the wavelength modes $\{\lambda_\mathbf{k}\}$.

The well-defined positive frequency subspace and mode functions given by Eq. (15) permit a time-translation operator to be defined via the Hamiltonian operator $\hat{H} = i\hbar\partial/\partial t$, with time parameter t [14,21,22].[20] Given these wavelength limitations and using the time translation operator, the field $\varphi(\mathbf{x}, t)$ induces a propagation from $\Sigma_0 \to \beta^+$ defined by the Feynman propagator for a virtual exchange particle as [22]

$$\langle 0| e^{-t_1 i\hbar\partial/\partial t}\varphi(\mathbf{y}_+)e^{-(t_1-t_0)i\hbar\partial/\partial t}\varphi(\mathbf{x}_{eh})e^{-t_0 i\hbar\partial/\partial t}|0\rangle \equiv i\triangle_F^+(y_+/x_{eh}), \tag{16}$$

which represents a virtual exchange field fluctuation created at spacetime event $(\mathbf{x}_{eh}, t_0)$ and propagated to $(\mathbf{y}_+, t_1)$ where it is annihilated. Thus when the positive frequency subspace is well-defined in a spacetime region, we can use the vacuum ket $|0\rangle \equiv |0_M\rangle$ and write Eq. (16) as representing a propagation of modes with $\lambda_\mathbf{r} \ll \rho_{eh}$ - ***if such a propagation were a physically viable phenomenon***. This leads to the question as to whether a phenomenon such as that represented by Eq. (16) can in fact exist from $\Sigma_0 \to \beta^+$.

c) *Proposition 3*

> **Necessarily spacelike paths exist which originate on a black hole's event horizon and extend to its exterior.**

This proposition follows from recognition that Σ_0 is a hypersurface composed of the null geodesics of null cones $\{\mathsf{N}\}$ tangent to Σ_0 [2,5-6,9,21], and that spacetime exists seamlessly across Σ_0 from β^- to β^+ [2-3,9]. If a path $\mathfrak{C}$ originating on Σ_0 were to proceed to β^-, then $\mathfrak{C}$ could be timelike, lightlike or spacelike. $\mathfrak{C}$ could pass through the interior $\mathfrak{N}$ of the null cone N with $\mathfrak{N} \subset \beta^-$, or alternatively pass from the origin $\mathfrak{O}$ of N with $\mathfrak{O} \in \Sigma_0$ while passing in a region $\mathfrak{R}$ with: $\{\mathfrak{R} \not\subset \mathfrak{N}$ but $\mathfrak{R} \subset \beta^-\}$. $\mathfrak{C}$ could also pass on any lightlike path emanating into β^-.

If the path originating on the event horizon remains on the event horizon, then it is a null path [2,6]. If however $\mathfrak{C} \to \beta^+$ from $x_{eh} \in \Sigma_0$, then it must extend to some exterior region $\mathfrak{J}$ of the null cone N originating at $x_{eh} \in \Sigma_0$ with null geodesic tangent to Σ_0. This follows because $\mathfrak{C} \to \beta^+$ while $\Sigma_0 \not\subset \beta^+$ and $\mathfrak{N} \not\subset \beta^+$[21]. Thus $\mathfrak{C} \to \beta^+ \subset \mathfrak{J} \not\subset \mathfrak{N}$, and it follows that $\mathfrak{C}$ must be spacelike.

[19] Ref. [4], Sec. 6.5.2.
[20] Ref. [14], Box 2-3 & Sec 3.1.5; Ref. [21], Sec. 30.4.
[21] $\mathfrak{N} \cap \beta^+ = \emptyset$.

An alternative way of conceiving of this proposition is to note that the spacetime curvature at Σ_0 warps all timelike world lines into β^- [2,3,9]. Since only null (lightlike) paths make up the event horizon, the only paths remaining that could proceed from $\Sigma_0 \to \beta^+$ must be spacelike in character. Since spacetime is extant across the event horizon of a black hole [2,3,5,9], spacetime paths exist there and so the paths proceeding from $\Sigma_0 \to \beta^+$ must be spacelike in character. Since they are spacelike in character, it further follows that they must represent paths for velocities greater than the speed of light of anything professing to travel from $\Sigma_0 \to \beta^+$ [2,3,27].

d) Proposition 4

> **The special theory of relativity's speed of light constraint does not apply to virtual exchange particles.**

The special theory of relativity sets the speed of light *in vacuo* as the ultimate speed [20,21,26,27,33].[22] However, quantum field theory requires that some set of modes of a virtual exchange particle travel faster than the speed of light between any two events [14,24].[23] This phenomenon is made evident when considering the integration of particle interaction transition amplitudes. These amplitudes are integrated over all spacetime [13-15], and include propagations of the virtual exchange particle outside the light cone of the source's spacetime locus.[24]

Virtual exchange particles exist off their own mass shell. It follows that virtual exchange particle modes can have arbitrarily high momenta for any spacetime propagation.[25] There then must exist some set of virtual exchange particle modes which when propagating from a source propagate on spacelike paths off the null cone emanating from the source's spacetime locus.[26] Thus to hold that virtual exchange particles cannot travel faster than the speed of light is in contradistinction to the basic tenets of quantum field theory.[27]

When the virtual exchange particle's momentum $\mathbf{p} = \hbar\mathbf{k} = h/\lambda_{\mathbf{k}}$ increases in the superposition of field modes, each mode's angular frequency $\omega_{\mathbf{k}}$ simultaneously increases in order to maintain the condition[28]

$$\omega_{\mathbf{k}} = (\mathbf{k}^2 + m_Q^2)^{1/2}. \tag{17}$$

Thus each successive field mode $\varphi_{\mathbf{k}}(x,t)$ has an increased momentum corresponding to a smaller wavelength $\lambda_{\mathbf{k}}$. In turn, the smaller $\lambda_{\mathbf{k}}$ corresponds to a greater speed - denoted as $d_{\lambda_{\mathbf{k}}}$- for the associated field mode. This notion is analogous to the quantization of harmonic oscillator eigenfunctions - each successive eigenfunction oscillating more rapidly, corresponding to an increase in momentum of the system and so an increase in the speed associated with each successive eigenfunction.[29]

Another way to conceive of this phenomenon is to consider the Heisenberg uncertainty relation's applicability [14,21-23].[30] Virtual exchange particles with off-shell mass m_Q may exist off their own mass-shell (0 in our case) so long as they abide by the duration-constraint [22]

[22] See Ref. [33], Sec. 2.3 for discussion; Also Ref. [21], Sec. 17.7-9.
[23] Ref. [24], pp. 94-6 & Fig. 61.
[24] Ref. [15], Ch. 1.4; Ref. [14], Sec. 8.6.
[25] Ref. [16], Ch. 14, "Asymptotically Free Partons"; Ref. [21], Sec. 26.6.
[26] See Ref. [21], Sec. 26.6 & 18.7, for such a depiction.
[27] Consider Ref. [23], p. 148, in which virtual exchange particles are considered to be limited by the speed of light. See also Ref. [22], p. 160, ambiguously noting a "finite velocity" for virtual exchange particles.
[28] Sec. II.iv, herein.
[29] Ref. [19], Sec. 10.7.
[30] See Ref. [22], Sec. 17.4 for a discussion.

$$\triangle t \lesssim \hbar / E_Q. \tag{18}$$

Thus $\mathbf{p} = \hbar\mathbf{k}$ may take on any value so long as the constraints of Eq. (17) {equivalently Eq. (7)} and Eq. (18) are met.

IV. Writing $i\triangle_F^+(y_+/x_{eh})$

Given the above propositions, we are nearly in position to write the Feynman virtual exchange particle propagator for the propagation of a virtual exchange particle from a spacetime event x_{eh} on the event horizon to a spacetime event y_+ in the exterior of the black hole, which can then be extended to $i\triangle_F^+(y_+/x_-)$.

To effect this, a curvature calibration procedure and a quantum-classical interface are appended to the technology, respectively called: 1) radius of curvature calibration (ρ-calibration), and 2) quantum field truncation (qf-truncation).

a) ρ-calibration procedure

We can conceive of the possibility of some set of virtual wavelengths $\{\lambda_\mathbf{r}\}$ in the mode expansion of our scalar field at Σ_0 having the elemental properties: $d_{\lambda_\mathbf{r}} > c$ and $\lambda_\mathbf{r} \not\ll \rho_{eh}$. That is, they will have escape velocities greater than the speed of light at the event horizon yet have wavelengths insufficiently small when compared to the radius of curvature there. In this scenario Proposition (2) would prevent use of a unique vacuum state $|0\rangle$.

However, since at any non-zero r-coordinate z there will exist some set $\{\lambda_\mathbf{j}\}$ of wavelengths such that $\forall \lambda_\mathbf{j} \in \{\lambda_\mathbf{j}\}$, $\lambda_\mathbf{j} \ll \rho_z$,[31] we can calibrate (adjust) ρ_{eh} so that the set $\{\lambda_\mathbf{k}\}$ of wavelengths in $\varphi(\mathbf{x}, t)$'s mode expansion which has the elemental property that $\forall \lambda_\mathbf{k} \in \{\lambda_\mathbf{k}\}$, $d_{\lambda_\mathbf{k}} \geq c$ will also have the property that $\lambda_\mathbf{k} \ll \rho_{eh}$.

We calibrate ρ_{eh} accordingly so as to refine definition of the set $\{\lambda_\mathbf{k}\}$:

Box I - the set $\{\lambda_\mathbf{k}\}$ defined
$\{\lambda_\mathbf{k}\}$ is the set of all $\lambda_\mathbf{k} = 2\pi/\mathbf{k}$ of $\varphi(\mathbf{x}, t)$ with set properties: $\lambda_\mathbf{k} \ll \rho_{eh}$ and $d_{\sup\{\lambda_\mathbf{k}\}} = c$.

In this definition $\sup\{\lambda_\mathbf{k}\}$ is the maximum wavelength of the set $\{\lambda_\mathbf{k}\}$. Since $d_{\lambda_\mathbf{k}} \to \infty$ as $\lambda_\mathbf{k} \to 0$,[32] all $\lambda_\mathbf{j} \in \{\lambda_\mathbf{k}\} \neq \sup\{\lambda_\mathbf{k}\}$ will have $d_{\lambda_\mathbf{j}} > c$.

b) Qf-truncation

Quantum field horizon truncation is conjectured to occur when a virtual exchange particle is created by a source at Σ_0 and interacts with the black hole there. This interaction results in the splitting or "filtering" of the superposition of momenta/wavelengths constituting the virtual vacuum field fluctuation (virtual exchange particle), which are given in the quantum field mode expansion of Eq. (1).

Qf-truncation is a distinctly quantum phenomenon the likes of which occurs pervasively throughout the quantum world, thereby providing support for the conjecture of qf-truncation at the event horizon of a black hole. To list a few examples: 1) the interaction of a quantum wavefunction ψ with a potential barrier;[33] 2) the interaction of a photon quantum wavefunction with a beam-splitter;[34] 3) the Stern-Gerlach effect; 4) the splitting and filtering of a fermionic wavefunction by passing it through a thin sheet of material; 5) the

[31] Since ρ_z will be finite and smoothly change as z changes.

[32] See Sec. II.iv & III.iv, herein.

[33] Ref. [4], p. 77.

[34] Ref. [21], Sec. 21.7.

Bohm-Aharonov effect; 6) quantum scattering theory.[35]

The common characteristic in these examples is the interaction between a quantum and classical object, e.g.:

$$\psi \longleftrightarrow V(x) \text{ (Potential barrier)},$$

$$\psi \longleftrightarrow \mathbf{B} \text{ (Stern-Gerlach)},$$

$$\psi \longleftrightarrow \mathbf{A} \text{ (Bohm-Aharanov)}.$$

Likewise with qf-truncation we have the interaction between a quantum object (quantum field) and a classical object (black hole):

$$\varphi(\mathbf{x}, t) \longleftrightarrow R_{\mu\nu\alpha\beta}.$$

Propositions (1)-(4) play a critical part in manifesting qf-truncation. Proposition (1) requires an interaction between the virtual exchange particle and the black hole. Proposition (2) sets the stage for conditioned representation of the virtual exchange particle mode propagation. Propositions (3) & (4) permit a limited propagation to occur off the event horizon and into β^+.

Interaction between a virtual exchange particle emitted at event x_{eh} and the black hole's gravitational field at Σ_0 splits the virtual quantum field disturbance into a part which is trapped by the black hole and a part which propagates off the event horizon to some event $y_+ \in \beta^+$ where it is absorbed. We will designate the truncated portion of $\varphi(\mathbf{x}, t)$ representing the $\mathbf{k}$-modes propagating on the paths $\Sigma_0 \to \beta^+$ as $\varphi_\tau(\mathbf{x}, t)$.

To obtain a definition of $\varphi_\tau(\mathbf{x}, t)$ we revisit Eq. (1)'s range of integration and the de Broglie relations of Eq. (2). The set of wave vectors $\mathbf{k}$ making up the superposition of $\varphi_\tau(\mathbf{x}, t)$ are related to the set $\{\lambda_\mathbf{k}\}$ of **Box I** in the following set Z to be integrated over in forming $\varphi_\tau(\mathbf{x}, t)$:

$$\mathsf{Z} = \left[\mathbf{k} : {}^{2\pi}/_{\mathbf{k}} = \lambda_\mathbf{k} \in \{\lambda_\mathbf{k}\}\right], \tag{19}$$

with which we can define the mode expansion of the truncated quantum operator field $\varphi_\tau(\mathbf{x}, t)$ as

$$\varphi_\tau(\mathbf{x}, t) = \int\limits_{\mathsf{Z}} \frac{d^3\mathbf{k}}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\omega_\mathbf{k}}} [e^{-ikx} \hat{a}^-_\mathbf{k} + e^{ikx} \hat{a}^+_\mathbf{k}]. \tag{20}$$

With this definition, ρ-calibration is incorporated into qf-truncation via the range of integration. This residual portion of the truncated quantum field at the event horizon necessarily exists, and creates a set of virtual exchange modes on Σ_0 which can propagate to β^+. Therefore $i\triangle^+_F(y_+/x_{eh})$ will constitute a non-zero amplitude which can be written as

$$i\triangle^+_F(y_+/x_{eh}) \; \equiv \; \langle 0| \, \varphi_\tau(y_+)\varphi_\tau(x_{eh}) \, |0\rangle > 0. \tag{21}$$

This is the two-way membrane hypothesis,[36] which predicts that there is a residual mode propagation of the virtual exchange fluctuation from the black hole's event horizon to its exterior. This hypothesis thus necessarily entails viewing the event horizon of a black hole as a limited two-way membrane, vice its traditional conception as a one-way membrane from which nothing can escape.

[35] Ref. [18], Sec. 19; Sec. 23; Sec. 24 & Ch. 7.
[36] Considering Eq. (3), it is seen that qf-truncation significantly reduces the amplitude for the occurrence of the subject interaction between the source at x_{eh} and sink at y_+.

c) Extension to $\mathbf{x}_-$ *and virtual cosmic censorship*

Extending Eq. (21) to a non-singular spacetime event $x_- \in \beta^-$ is straightforward. One redefines the sets $\{\lambda_\mathbf{k}\}$, Z and the supremum of $\{\lambda_\mathbf{k}\}$ so as to permit a more limited mode expansion of the truncated operator field, creating what is in effect a virtual event horizon with radius given by $h_{x_-} = 2GM/d^2_{\sup\{\lambda_\mathbf{k}\}}$, resulting in the equation $i\triangle_F^+(y_+/x_-) > 0$.

This extension reveals another phenomenon. At separations $h < R_s$ the escape velocity $s_{x_-} > c$.[37] We further have that

$$\begin{gathered} s_{x_-} \to \infty \quad as \quad h \to 0, \\ \rho \to 0 \quad as \quad h \to 0, \\ \{\lambda_\mathbf{k}\} \to \emptyset \quad as \quad h \to 0. \end{gathered} \tag{22}$$

From the relations of Eq. (22) it immediately follows that

$$\lim_{h \to 0} i\triangle_F^+(y_+/x_-) = 0. \tag{23}$$

This is the virtual cosmic censorship hypothesis. As x_- draws closer to the singularity, a larger set of virtual particle modes is trapped. Per Eq. (23), in the limit of the singularity all virtual exchange vacuum field fluctuations are trapped, and thus none escape to β^+ in this limit. Since no virtual exchange modes can escape in this limit, there can be no conceivable exchange propagation off the black hole's singularity. Since all measurements occur via the exchange of virtual exchange particles, it follows that the singularity is hidden from the view of any observer external to it.

V. Conclusion

This paper's straightforward propositions generate important results. Given the paper's technology it is deduced that a black hole's event horizon is not a one-way membrane, but a limited two-way membrane in which there exists the real possibility for a residual transmission off the event horizon (as well as from the black hole's interior) of a subset of the wavelength modes which comprise a virtual vacuum field exchange fluctuation. In addition, if we maintain the quantum field theoretic presumption that the only feasible scientific method of observation is via the mechanism of virtual exchange particles, then we require a black hole's singularity to be hidden from direct observation.

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[37] Ref. [12], Sec. 13.5 (p. 508).

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