Time symmetry and measurement at the event horizon of a black hole.

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Abstract:
In this initial investigation, we establish rules for the measurement of Lorentz invariant potential at the event horizon of a black hole. We define "measurement" as the application of real boundaries to complex geometric potentials. Measurement in this context does not require observers and, therefore, measurement boundaries can be applied to both objects in real space and geometric potentials in the complex plane. The separation of quantum and wave based time operators at the horizon allows us to model the interior of a black hole as an asymptotic quantum well for Lorentz invariant potential, based on least-time principles and the collapse of the wave function. As we will demonstrate, defining the boundary between quantum and relative time operators also requires a separation of charged and uncharged potentials at the transition boundary. We discuss the repercussions resulting from the discarding of charged potential, at the horizon, as backward transition radiation (BTR). We then tie (BTR) to a mechanism for synchrotron emission and x-ray generation in the inner accretion disk. We also connect (BTR) to AGN heating and cooling feedback loop models for jet generation at the poles. We finish with a discussion regarding uncharged quantum binary potential and the polarity and tensor-driven nature of gravitational waves. We connect this model to observations of from The William Herschel telescope and the Chandra, NUSTAR and LIGO observatories.

Keywords: Black hole Horizon; Lorentz invariance; Black hole model; Quantum gravity; Statistical mechanics; White hole Dynamics; Black hole reflection; Charged accretion disk; Gravitational wave polarity; Quantum black hole

1.00 - Introduction

Many black hole models have explored the limits at event horizon and come to the conclusion that the interior source for gravitational potential may be quantum in nature (Bekenstein, 1975)(Hawking, 1974)(Hogan, 2015)('t Hooft, 2016)(Rovelli, 2019)(Dray and Rovelli, 2019). In this letter, we model the dynamics at the event horizon of a black hole using rules for the measurement of Lorentz invariant potential at the collapse of the wave function. Rules for measurement set a hard line between quantum boundaries with *relative* potential and relative boundaries with *quantum* potential. Utilizing the separation of Markovian and Bayesian time operators allows us to define the boundary between complex potential and relative potential at the horizon of a black hole (BH). We constrain all measurements to the same boundaries established by general relativity that constrain Lorentz invariant quantities. In this model, fundamental quantum states of potential are determined by measurement, establishing their probability in in real space. We define "measurement" as the application of real boundaries to complex geometric potentials. Measurement in this context does not require observers and, therefore, measurement boundaries can be applied to both objects in real space and geometric potentials in the complex plane. If we consider the event horizon as the "measurement" that determines this boundary; we can employ rules for measurement to define the line between quantum and relative time operators at the horizon. In this model, least time boundaries are defined by the distinct differences in the way we measure quantum or relative time operations. Bayesian time is the time of Einstein and is held to all the boundaries of Minkowski space and special relativity (Einstein, 1905), Minkowski, 1908). It flows based on the concept of cause and effect. All wave functions operate using this time operator and are bounded by the speed of light. This includes all the Bob and Alice scenarios used to describe the relative measurements taken by different observers. Markovian time operators have very different boundaries and are not limited by the speed of light. Markovian time operators are quantum in nature and carry specific rules for time transition operators   Entanglement is best described using Markovian time operators (Wheeler,b 1978)(Nambu, 1973)(Meyn and Tweedie 1993). These are the basic rules for the measurement and interaction of Markovian and Bayesian time operators that we will employ for this discussion:

- *Two or more simultaneous measurements constitute a single relative measurement*[1]
- *Quantum measurements require a quantum time operator*
- *Quantum measurements cannot be simultaneous*
- *Simultaneous measurements that combine two or more different quantum time operators must use a shared relative time operator*.

[1]*(Einstein places limits on simultaneous measurement based on the limits of relative observation. In this model, all simultaneous measurements must establish a common relative time operator. This is essentially a re-statement of the same principle. We are bounded by the same limits established by Lorentz invariance and Minkowski space. This is discussed further in Sections-1.01 and 1.02)*

These simple rules will help us to define boundary conditions for unmeasured vector potential which establish the boundaries for measured gravitational potential. General and special relativity tell us that there are no simultaneous measurements and the relative velocity of observers limits the context of the measurement. Our first rule for measurement seems to conflict directly with this basic limit to any simultaneous measurement imposed by general relativity. In fact, our first rule states that any simultaneous measurements must take into account the relative time operators of each measurement. This is actually in agreement with Einstein's principles regarding simultaneous measurement. In effect, we create quanta that respond to measurement to determine quantum states. Quantum measurements simply require a quantum time operator. This single requirement serves to separate quantum and relative time operations and implies the need for a discrete Markov chain in balance with a Bayesian wave function. We then draw a hard line between these two time functions and the center of least time through the use of a Nambu-Poisson symmetry and phased Hamiltonians representing the center of least time for measured potential. As we will attempt to demonstrate; in a system determined by measurement there are no zero states or infinities to renormalize because *both zero and infinity are not measurable states*. *(See: Addendum: A for an extended discussion of the minimums for measurement in real space)*. Setting minimums for measurement allows us to create a symmetry between quantum and relative potential based on least-time boundaries (Bekenstein, 1975)(Hawking, 1974)(Hogan, 2015)(t'Hooft, 2016)(Rovelli, 2019). In this investigation we apply the principle of minimum time and action to help us define the boundary between Markovian and Bayesian operators at the horizon. We define a time symmetry between quantum and wave-based time operators employing the same Nambu-Poisson symmetry that is the basis of electroweak renormalization (Nambu, 1973)(Takhtajan, 1994). Generalized Nambu-Hamilton equations of motion involve two Hamiltonians and an evolutionary time operator. They establish how the system evolves using the principle of least time represented by the Hamiltonians within the equation:

$$\frac{df}{dt} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_2 \to t_1)\} \tag{01}$$

Where, $f(t_2 \to t_1)$ represents the least-time evolution operator of the two symmetric Hamiltonians.

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_1)\ \} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_2)\} \tag{02}$$

We establish Markovian boundaries to the quantum (left) side of the Nambu-Poisson symmetry by equating an identity matrix to a time operator. Markovian transition matrices allow us to create boundaries for least-time potential based on the restrictions of the matrix. One of the primary advantages of working with *unmeasured* potential is that it allows us to maintain positive boundaries for *measured* potentials in Markovian transition matrices. Unmeasured potential can also placed in the complex plane as an asymptotic boundary condition for measured potential. We can represent the boundaries for unmeasured potential using the set of complex numbers and Hermitian matrices. We start with a transition matrix using real complex numbers representing the opposite potentials needed to reflect conservation boundaries:

$$\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \tag{03}$$

This transition matrix establishes boundary conditions for all measured binary potential. We can associate this basic binary requirement to uncharged Markovian transition probability matrices:

$$f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} = f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}^2 = f\begin{bmatrix} +1 & -1 \\ -1 & +1 \end{bmatrix} = f(t_n) \tag{04}$$

Where the value $f\begin{bmatrix} +1 & -1 \\ -1 & +1 \end{bmatrix}$ represents the least-time boundaries for unmeasured Markovian potential and $f(t_n)$ is the quantum evolutionary time operator. In our separation of quantum and relative operations, required by our rules for measurement, any Markovian binary boundary condition requires this basic identity as a time operator:

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\} = -\left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\right\} \tag{05}$$

Using this matrix as the Markovian transition time operator allows us to create asymptotic boundaries for unmeasured potential that will set the positive boundaries for measured potential. In our model, quantum time operators, act as uncharged, and asymptotic, binary boundaries for Bayesian charge and wave-based potentials. Charged potentials combine separate quantum measurements into a single measurement and, therefore, based on our rules, we must use a Bayesian wave function as the time operator for any charged potential. This means that all charges, positive, negative and neutral must be contained on the Bayesian side of the NP symmetry. We can show this best by returning to Markovian transition matrices representing charged and wave-based potential. If we include all three charges of the standard model on the Bayesian side of our symmetry we create a transition identity for the wave packet. We can now use the set of complex numbers to represent neutral charge potential:

$$t_\psi = \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix} = \begin{bmatrix} 1^+ & i & 1^- \\ i & i & i \\ 1^+ & i & 1^- \end{bmatrix} \tag{07}$$

In our model we use $(1^+, 1^-, 1)$ to represent positive, negative and neutral currents. We use a quantum fundamental state as the carrier for neutral currents, because it carries electrodynamic unmeasured potential without the boundaries set by Coulomb rules and charges. Vector potential works as a quantum operator. By assigning the appropriate quantum or wave-based time operations we can tie these equations to our symmetry between time operators. If we examine the diagonal representing this time operator as a position vector (r) then the wave function represents the boundary conditions for the general Schrodinger equation:

$$\hbar i \frac{\partial}{\partial t} |\psi(r,t) \;\; = \;\; \mathcal{H}|\psi(r,t) \tag{07}$$

$$\hbar i \frac{\partial}{\partial t} |\psi(r,t) \;\; = \;\; \mathcal{H}|t_{\psi} \tag{08}$$

$$\mathcal{H}|t_{\psi} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_{\psi})\} = -\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\ \} \tag{09}$$

$$-\left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\ \right\} = \left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix}\right\} = i\frac{\partial}{\partial t}|\psi(r,t) = \mathcal{H}|\psi(r,t) \tag{10}$$

To define the boundaries that are set by our rules for measurement, we begin by establishing a boundary between least-time quantum and wave-based time operators in a Nambu-Poisson symmetry. We use an integer-based Markovian operator as our quantum time operator and a Bayesian wave packet as the relative time operator, :

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\ \} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_{\psi})\} \tag{11}$$

Where the value (t_n) represents a discrete Markovian time operator and (t_{ψ}) represents a wave-based Bayesian time operator. Adding sin and cos values to each Hamiltonian will allow us to translate between wave and geometric quantum operators in real space (See: *Addendum A*):

$$-\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_n)\ \} = \{\mathcal{H}_{sin}\, , \mathcal{H}_{cos}, f(t_{\psi})\} \tag{12}$$

We can assign quantum or relative time operators to actions both in and outside the collapse of the wave function and map them as Lorentz-invariant Hamiltonians using quantum and relative time operators and the Green's functions:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_n)\ \} = \nabla^2 G\{\mathcal{H}_{sin}\, , \mathcal{H}_{cos}, f(t_{\psi})\} \tag{13}$$

In the next section we will explore how we can use these principles to establish the quantum gravitational potential at the horizon.

1.02 - Method

In this model, gravity is positioned as a quantum asymptotic boundary condition for classic relative potential based on the principle of least time. This seems like a departure from general relativity, but it actually is just a rephrasing of the boundaries without the need for fields. We establish the same asymptotic boundaries established by Lorentz invariance using the stress tensor and the gravitational constant as least-time boundaries. In our model, gravimetric *measured* potential is defined using the same boundaries established by Einstein and Minkowski (Einstein, 1905), Minkowski, 1908). In fact, if we think of classic potential in its most recognizable form we can turn to Einstein's famous equivalency of energy to mass:

$$E = mc^2 \tag{14}$$

This equivalency of energy to mass is basically a statement that equates energy to potential. Our use of NP symmetries allows us to translate Lorentz invariant potential between quantum and relative states. In effect, we treat Lorentz invariance as a quantity that we can translate between quantum and relative time operations using the hard line set by our rules for measurement. This allows us to fulfill the unique requirements of special relativity as well as working outside the collapse of the wave function and with quantum boundaries. Let's return to our original equivalency of charged potential and uncharged binary boundary conditions. In this instance we will measure the asymptotic boundary for classical potential momentum based on a quantum center of least time:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \leftarrow \nabla^2 G f_\psi \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix} \tag{15}$$

While relative gravitational potential, and the relative time operator, are always limited by the wave function and the speed of light, quantum gravitational potential is not. We define this quantum minimum for gravitational potential as an uncharged asymptotic boundary for classic potential. Lorentz invariant quantities define the minimum measurement for any classic potential. By attaching vectors to Lorentz invariant classic potential, we can portray the center for least time as an uncharged quantum binary boundary condition for relative potential based on conservation of momentum:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & mc^2 \\ mc^2 & 1 \end{bmatrix} \leftarrow \nabla^2 G f_\psi \begin{bmatrix} 1^- & mc^2 & 1^+ \\ mc^2 & mc^2 & mc^2 \\ 1^- & mc^2 & 1^+ \end{bmatrix} \tag{16}$$

Beginning with our description of classic momentum, we can use our NP symmetry, and the gravitational constant, to define a center for least time for gravitational statistical potential. We can use the Einstein's momentum and energy stress tensor to construct a minimum for quantum measurement of these momentum vectors and boundaries. Einstein uses the energy stress tensor ($T_{\mu\nu}$) to represent the asymptotic source of gravitational potential. Because it represents the center of least time for any Lorentz-invariant momentum, we can use the energy-stress tensor as our quantum time operator for vectored momentum. We can also use it on the charged side of our symmetry, but we change to the opposite vector constrained by the binary uncharged boundary condition:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & \overleftarrow{T_{\mu\nu}} \\ \overleftarrow{T_{\mu\nu}} & 1 \end{bmatrix} = \nabla^2 G f_\psi \begin{bmatrix} 1^- & \overrightarrow{T_{\mu\nu}} & 1^+ \\ \overrightarrow{T_{\mu\nu}} & \overrightarrow{T_{\mu\nu}} & \overrightarrow{T_{\mu\nu}} \\ 1^- & \overrightarrow{T_{\mu\nu}} & 1^+ \end{bmatrix} \tag{17}$$

Returning to our original NP symmetry:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(\overleftarrow{T_{\mu\nu}})\} = \nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi(\overrightarrow{T_{\mu\nu}})\} \tag{18}$$

In our model we represent gravity as an uncharged boundary condition for vectored classic momentum. To set a hard line between quantum and relative time operators, let us return to the use of Einstein's gravitational constant[7]:

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\left(\overleftarrow{T_{\mu\nu}}\right)\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{8\pi G_g}{c^4}\left(\overrightarrow{T_{\mu\nu}}\right)\right)\right\} \tag{19}$$

We can use the gravitational coupling constant (α_G) to help us move all charge to one side of our equation:

$$G_g = \frac{\alpha_G hc}{2\pi {m_e}^2} \tag{20}$$

Where (G_g) is the gravitational constant and α_G is the gravitational coupling constant:

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\, T_{\mu\nu}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{8\pi\alpha_G hc}{2\pi {m_e}^2 c^4}\, T_{\mu\nu}\right)\right\} \tag{21}$$

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\, T_{\mu\nu}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{4\alpha_G h}{{m_e}^2 c^3}\, T_{\mu\nu}\right)\right\} \tag{22}$$

Moving all light speed and charge boundaries to the Bayesian side of the equation:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(8\pi G_g\, T_{\mu\nu})\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{4\alpha_G hc}{{m_e}^2}\, T_{\mu\nu}\right)\right\} \tag{23}$$

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(2\pi G_g\, T_{\mu\nu})\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{\alpha_G hc}{{m_e}^2}\, T_{\mu\nu}\right)\right\} \tag{24}$$

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{\alpha_G}{2\pi}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{G_g {m_e}^2}{hc}\right)\right\} \tag{25}$$

Adding back the stress tensor allows us to establish a geometric center of least time and gravitational potential that also separates quantum Markovian and wave-based Bayesian boundaries and uncharged and charged potentials:

$$f_n\left(\frac{\alpha_G}{2\pi}\right) T_{\mu\nu} \;\leftarrow f_\psi\left(\frac{G_g {m_e}^2}{hc}\right) T_{\mu\nu} \tag{26}$$

Any particle with mass approaches the horizon of as BH as a relative particle operating in relative space because it carries charge. According to our rules for measurement, *when any relative potential crosses the horizon it is entering an area that is outside the collapse of the wave function. It must therefore discard any Bayesian, and relative, time operators*. It is this separation of uncharged and charged potentials that we use to define measurement minimums at the horizon. As we know from basic electrodynamics, when any charged potential crosses a boundary with two different dielectric constants it discards charge information as transition radiation (Ginzburg, 1982). Our rules for measurement require that any a charged potential crossing the must leave behind relative charge information. In our model, charged particles must discard charge information outside of the horizon as entangled momentum; then they can cross the horizon as binary uncharged potential. We can relate this directly to relativistic and non-relativistic electrons crossing a reflective boundary. For the case of transition radiation involving an electron with velocity $(-v)$ reflecting from an ideal with a reflective angle θ, we can state:

$$\mathcal{W}_1(\omega, \theta) = \frac{e^2 v^2 sin^2\theta}{\pi^2 c^3 \{1-(v/c)^2 cos^2\theta\}^2} \tag{27}$$

$$\mathcal{W}_1(\omega) = \frac{e^2}{\pi c}\left\{\frac{1+(v/c)^2}{2v/c} ln \frac{1+v/c}{1-v/c} - 1\right\} \tag{28}$$

In the non-relativistic case where $(v \ll c)$:

$$\mathcal{W}_1(\omega, \theta) = \frac{e^2 v^2 sin^2\theta}{\pi^2 c^3} \tag{29}$$

$$\mathcal{W}_1(\omega) = \frac{4e^2 v^2}{3\pi c^3} \tag{30}$$

In an the ultra relativistic limits that we find at the event horizon where $(v \rightarrow c)$ we get:

$$\mathcal{W}_1(\omega) = \frac{e^3}{\pi c} ln \frac{2}{(1-v/c)} = 2\frac{e^2}{\pi c} ln \frac{2E}{mc^2} \tag{31}$$

Therefore, the reflective backward transition radiation can be stated as having an energy:

$$E = \frac{mc^2}{\sqrt{[1-(v^2/c^2]}} \gg mc^2 \tag{32}$$

Backward transition radiation is also proportional to the Lorentz factor, (L_f) and the transition wavelength (λ) for any mass crossing a boundary in relative space where $E \gg mc^2$:

$$L_f = 2\lambda \left(\frac{E}{mc^2}\right)^2 \tag{33}$$

$$\gamma = 2\lambda \left(\frac{E}{mc^2}\right)^2 \tag{34}$$

In our model, the interior of the black hole becomes an uncharged binary boundary condition defined by a single quantum time operator and bounded by relative space, therefore:

$$\frac{\gamma}{2} = \lambda \left(\frac{E}{mc^2}\right)^2 \tag{35}$$

When we add the consideration that the discarding of charge at the boundary is relativistic. We can tie charge BTR and Lorentz factor to the generation of synchrotron radiation based on the Larmor formula for relativistic charged point particles travelling in a curved path. To simplify the discussion, let us begin with the non-relativistic Larmor formula for emitted radiation from a charged particle in a magnetic field (Larmor, 1897) :

$$P_{rad} = -\frac{dE}{dt} = \frac{2}{3}\frac{q^2 a^2}{c^3} \qquad (36)$$

Where the radiated power $= -\frac{dE}{dt}$ generated by an accelerated charge $= q$ is proportional to the square of the acceleration $= a$. To simplify our discussion of non-relativistic and relativistic particles we can state the covariant form using the Lorentz factor $= \gamma$. This will allow us to work with Lorentz invariant quantities and four-momentum:

$$dP_\lambda = \frac{2q^2}{3c^3}\frac{1}{(mc)^2}\frac{dp_\mu}{d_\tau}\frac{dp_\mu}{d_\tau}dx_\lambda \qquad (37)$$

Here the energy momentum four-vector $dp_\mu = (ic, \mathbf{p})$ and d_τ is the invariant proper time such that $d_\tau = \gamma^{-1}dt$. In our model we use Lorentz invariant potentials to describe all phases of energy. Just like the transition between two different dielectric mediums, charge four-momentum is discarded as any particle crosses into uncharged quantum space. The result is the escape of synchrotron radiation and the resulting generation of x-rays in the inner accretion disk (Kumar et al., 2014)(Kumar and Chattopadhyay 2017). Discarded charge potential is conserved as backwards ultra relativistic transition radiation entangled with the interior binary potential. Unlike the very small effect of Hawking radiation (S. Hawking, 1974) our model shows how charged particles discard large amounts of entangled potential before crossing any boundary between binary quantum and relative space as synchrotron radiation.

1.03 - Observational evidence

Because of the NP symmetries we have attached to this model, the polarity of the uncharged quantum interior has a number of other repercussions in relative space which can be tied to direct observation. Chiefly among these attributes the definition a mechanism for the emission of synchrotron radiation and the resulting creation of bi-polar jets at the poles. Many models for exist that characterize bi-polar jets at the poles coming from a rotating accretion disk as opposed to being generated by the magnetic field coming from the interior. These models are well summarized in two thorough, and well referenced papers(Kumar et al., 2014)(Kumar and Chattopadhyay 2017). In these papers, they reached a number of important conclusions regarding the nature of bipolar-jets and their connection to the accretion disk:

- *Jet states are deeply linked with the accretion states.*(Kumar et al., 2014)
- *The entire accretion disc does not generate jets, only the inner region forms jets.*(Kumar et al., 2014)
- *Relativistic jets can be obtained if radiation energy, as opposed to thermal energy, is supplied to the accretion disk.* (Kumar and Chattopadhyay 2017)

If we view the interior of the black hole as an uncharged binary boundary at the collapse of the wave function that separates quantum and relative potential; then the discarded BTR charge could generate a tremendous amount of outflowing potential producing synchrotron radiation and generating the power for jets at the poles. An x-ray heating source in the inner accretion zone, nearest the horizon, also fits well with with many models for jet generation from the inner accretion disk(Kumar et al., 2014)(Kumar and Chattopadhyay 2017). Here we quickly outline some of the observational data that fits our BTR model. Obviously, a more rigorous approach to the data is required, but it is beyond the scope of this initial letter.

CHANDRA*:*

The generation of synchrotron radiation from discarded charge potential at the horizon fits well with descriptions of the radiative cooling process. Chandra observations also show a rapid heating and cooling in the accretion disk due to x-ray emission. (H. Tananbaum et al., 2014):

"*In simplified terms, the AGN feedback cycle can be summarized as follows:*

1. accreting gas falls towards a SMBH and is heated, converting gravitational potential energy to radiation including X-rays;

2. jets are launched from the SMBH (BH spin may play an important role here), re-heating radiating gas to prevent runaway cooling and pushing aside infalling gas;

3. gas supply is diminished, jets turn off, and the SMBH returns to inactive state;

4. accretion resumes and the cycle starts over." (H. Tananbaum et al., 2014)

Our model for the generation of synchrotron radiation from discarded charge potential at the horizon fits well with this description of the radiative cooling process in AGN. The discarded charge could generate a tremendous amount of outflowing potential which would add to the angular momentum of the accretion disk and produce synchrotron radiation. This is supported by observations of outflowing "winds" shown in the Chandra data(J. Miller et al. 2017):

"*The absorption is primarily from H-and He-like ions with atomic number Z ranging from 8 to 28, indicating that the ions are relatively near to the BH, while the blue shifts and column densities indicate the presence of a dense wind. Detailed modeling shows that the observed wind cannot be driven by radiation pressure or photoionization. Heating by magnetic turbulent viscosity in the accretion disk is a possible cause for the wind. The wind carries away orbital angular momentum, allowing a large fraction of the gas in the disk to spiral inward to the BH*." (J. Miller et al. 2017)

NUSTAR, William Herschel Telescope*:*

Our model can also be tied to recent observations of a phase lag between x-ray data from NUSTAR and optical data from the William Herschel Telescope. Observations of jet flares at V404 Cygni were shown to exhibit such a lag (Gandi et al. 2017). Dr. Gandi's team characterized this as:

"*a lag between relativistic Lorentz factors powering x-ray emission at the accretion zone and the resulting creation of jets at the poles*(Gandi et al. 2017).

"*If the jet is fed by the accretion flow, a time lag between the x-ray emitting accretion flow and the shock synchrotron radiation is expected, with the magnitude of this lag depending mostly on the characteristic time scale of velocity* *fluctuations driving the shocks. Lags of order* τ*~0.1s may naturally be obtained under the assumption that the power spectrum of the distribution of shock Lorentz factors is similar to that powering the x-ray emission.*"(Gandi et al. 2017)

As we demonstrated earlier, BTR is also proportional to the Lorentz factor, (L_f) and the transition wavelength (λ) for any mass crossing a boundary in relative space.

LIGO:
A quantum polar interior can also be used to model the polarity of gravitational waves. Our model at the horizon also creates the binary boundary conditions driven by the stress tensor that fit well with the observational data from LIGO regarding polarity of gravitational waves. Our use of the stress tensor to define the horizon is also well supported by observations by recent data from the LIGO observatory(B. P. Abbott et al. 2017):

"The addition of Virgo has allowed us to probe the polarization content of the signal for the first time; we find that the data strongly favor pure tensor polarization of gravitational waves, over pure scalar or pure vector polarizations."(B. P. Abbott et al. 2017)

1.04 - Discussion

By applying our rules for the measurement of Markovian and Bayesian boundaries for Lorentz invariant potential at the horizon of a black hole, we have been able to draw a hard line between quantum and relative time operations and boundary conditions. *We have shown that the reflection of charged BTR potential can model quantum gravitational potential, synchrotron radiation, x-ray emission, and observations of angular momentum and energy added to the accretion disk in the form of outflowing "winds*". We have used these principles to model many of the primary aspects of BH observations; including x-ray emission, heating and cooling feedback loops in AGN observed by the Chandra `observatory(H. Tananbaum et al.), radiative sources for BH winds(H. Tananbaum et al. 2014)(J. Miller et al. 2017) as well as a source for radiative jets at the poles (Kumar et al 2014)(Kumar and Chattopadhyay 2017). We have also shown that a bi-polar quantum interior can model gravitational wave data from the LIGO observatory, indicating the polarity of gravitational waves as tensor driven (B. P. Abbott et al. 2017).

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