The Hawking effect is short-lived in polymer quantization

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It is widely believed that the Hawking effect might hold clues to the possible, yet unknown, trans-Planckian physics. On the other hand, one could ask whether the effect itself might be altered by such trans-Planckian physics. We seek an answer to this question within a framework where matter field is quantized using polymer quantization, a canonical quantization technique employed in loop quantum gravity. We provide an exact derivation of the Hawking effect using canonical formulation by introducing a set of near-null coordinates which allows one to overcome the challenges posed by a Hamiltonian-based derivation of the Hawking effect. Subsequently, we show that in polymer quantization the Hawking effect is short-lived and it eventually disappears for an asymptotic future observer. Such an observer finds the duration of the Hawking effect to be few milliseconds for a solar mass black hole whereas it is few years for an ultra-massive black hole. Consequently, it provides a new way to resolve the so-called information loss paradox.

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Introduction.— The Hawking effect [1] continues to be an enigma in modern physics where an asymptotic future observer experiences a thermal emission, rather unexpectedly, emanating from a classical black hole. In statistical physics, the thermal emissions are known to arise from systems having large number of microscopic degrees of freedom. However, to describe black holes which are solutions of Einstein’s general relativity, only a few parameters are required. This perplexing property suggests that the Hawking effect might hold the key in understanding possible microstates of a black hole. These states are expected to arise from a possible, yet unknown, quantum theory of gravity and have been pursued extensively in different contexts [2].

It is well known that all prominent derivations of the Hawking effect rely on the properties of the trans-Planckian frequencies one way or other. On the other hand, it’s widely expected that our current understanding of trans-Planckian physics would need to be modified in order to tame the plaguing ultraviolet divergences. Therefore, one is led to ask whether the Hawking effect itself could survive these expected trans-Planckian modifications [3]. Besides, the evaporation of a black hole through Hawking radiation gives rise to the so-called information loss paradox [4, 5] which, according to the popular school of thought, threatens unitarity, a key pillar of quantum theory (see also [6]).

We seek an answer to the question within the framework of polymer quantization of matter field in the Schwarzschild geometry which is formed through a collapsing shell of matter. Polymer or loop quantization [7] is a canonical quantization technique which is employed in loop quantum gravity [8]. This quantization comes with a new length scale which would correspond to the Planck length in full quantum gravity. However, here we would employ this quantization only for the matter sector and treat the spacetime geometry as the classical entity [9], as done for the standard derivation of the Hawking effect. Similar studies in the context of the Unruh effect [10] has indicated significant modification [11, 12].

It turns out that there are major hurdles in pursuing a Hamiltonian-based derivation of the Hawking effect. The key reason behind these hurdles is the fact that thermal characteristic of the Hawking quanta is realized using the relation between the modes that leave past null infinity as ingoing null rays and the modes that arrive at future null infinity as outgoing null rays. Expectedly, the usage of the advanced and retarded null coordinates rather than the regular Schwarzschild coordinates, forms a crucial backbone for the standard derivation of the Hawking effect. However, these null coordinates do not lead to a true Hamiltonian that describes evolution of the relevant modes (see also [13]). In an earlier such attempt by Melnikov and Weinstein [14] who used Lemaître coordinates, the Hawking effect is understood indirectly through the property of the Green’s function. To the best of our knowledge an exact derivation of the thermal spectrum for Hawking radiation using Hamiltonian formulation is still lacking.

Schwarzschild spacetime.— Let us consider a Schwarzschild black hole which is formed after the collapse of a matter shell. The corresponding metric is

\begin{equation}
\text{ds}^2 = -\Omega dt^2 + \Omega^{-1} dr^2 + r^2 d\theta^2 + r^2 \sin \theta^2 d\phi^2,
\end{equation}

where \(\Omega = (1 - r_*/r)\). Here we have chosen natural units such that \(c = \hbar = 1\) and the Schwarzschild radius \(r_* = 2GM\). If one defines the so-called tortoise coordinate \(r_*\) such that \(dr_* = \Omega^{-1} dr\), then \(t - r\) plane of the Schwarzschild geometry becomes conformally flat. By a suitable choice of constant of integration, \(r_*\) can be written as \(r_* = r + r_0 \ln (r/r_* - 1)\). For later convenience, we define the advanced and retarded null coordinates \(v = t + r_*\) and \(u = t - r_*\) respectively.

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In addition to the collapsing shell of matter, we consider a minimally coupled, massless scalar field $\Phi(x)$ to represent the Hawking quanta $|1\rangle$, and which is governed by the action $S_\Phi = \int d^4x \left[-\frac{1}{2}\sqrt{-g}g^{\mu\nu}\nabla_\mu \Phi(x)\nabla_\nu \Phi(x)\right]$. For an observer at past null infinity $\mathscr{I}^-$, the scalar field operator can be expressed as

$$\hat{\Phi}(x) = \sum_{\omega} \left[f_\omega \hat{a}_\omega + f_\omega^* \hat{a}_\omega^\dagger\right], \tag{2}$$

where the set of ingoing field solutions $\{f_\omega\}$ forms a complete family on $\mathscr{I}^-$ along with the inner product $\langle -i/2 \int_S d\Sigma^a (f_\omega \nabla_a f_{\omega'} - f_{\omega'} \nabla_a f_\omega) = \delta_{\omega \omega'}$ where $\mathcal{S} = \mathscr{I}^-$. In order to render the inner product positive definite, only positive frequency modes, with respect to a canonical affine parameter along $\mathscr{I}^-$, are chosen. These modes can be written as

$$f_\omega(v) = \frac{1}{\sqrt{2\pi\omega}} r^{-1} e^{-i\omega v} Y_{lm}(\theta, \phi), \tag{3}$$

where $Y_{lm}(\theta, \phi)$ are spherical harmonics. The creation and annihilation operators are $\hat{a}_\omega$ and $\hat{a}_\omega^\dagger$, respectively. The vacuum state $|0\rangle$ is defined as $\hat{a}_\omega |0\rangle = 0$. Similarly for a future observer we can express $\hat{\Phi}(x)$ as

$$\hat{\Phi}(x) = \sum_{\omega} \left[p_\omega \hat{b}_\omega + p_\omega^* \hat{b}_\omega^\dagger\right] + \sum_{\omega} \left[q_\omega \hat{c}_\omega + q_\omega^* \hat{c}_\omega^\dagger\right], \tag{4}$$

where field solutions $p_\omega(u) = \frac{1}{\sqrt{2\pi\omega}} r^{-1} e^{-i\omega u} Y_{lm}(\theta, \phi)$ are purely outgoing and $(\hat{b}_\omega^\dagger, \hat{b}_\omega)$, $(\hat{c}_\omega^\dagger, \hat{c}_\omega)$ are creation and annihilation operator pairs at future null infinity $\mathscr{I}^+$ and event horizon respectively. The solutions $\{p_\omega\}$ have zero Cauchy data on event horizon whereas the solutions $\{q_\omega\}$ have zero Cauchy data on future null infinity $\mathscr{I}^+$. Hawking radiation. For derivation of Hawking effect, an essential relation between null coordinates on $\mathscr{I}^-$ and $\mathscr{I}^+$ with suitable choice of pivotal values (see FIG.1) is given by

$$-u = -v + 2r_s \ln(-v/2r_s). \tag{5}$$

The relation (5) crucially depends on the fact that there was no black hole when relevant ingoing modes departed from $\mathscr{I}^-$. For Hawking radiation, relevant modes originate from the region $|v| \ll 2r_s$ on $\mathscr{I}^-$ and for them the relation (5) can be approximated as

$$v \approx -2r_s e^{-u/2r_s}. \tag{6}$$

The Hawking effect is realized from the expectation value of number operator corresponding to the observer at future null infinity $\mathscr{I}^+$ in the vacuum state corresponding to the observer at past null infinity $\mathscr{I}^-$, and is given by

$$N_\omega \equiv \langle 0_-|\hat{b}_\omega^\dagger \hat{b}_\omega|0_-\rangle = \frac{1}{e^{2\pi\omega/\kappa} - 1}. \tag{7}$$

where $\kappa = 1/(2r_s)$ is the surface gravity at the horizon. This perceived phenomena of blackbody radiation at $\mathscr{I}^+$ is referred to as the Hawking effect with Hawking temperature $T_H = \pi/(2\kappa k_B) = 1/(8\pi GMk_B)$. We note that despite being mentioned frequently the derivation of the Hawking effect does not require any pair-production of particles nor these particles would have any Cauchy data on $\mathscr{I}^-$. Canonical formulation. In canonical formulation, field dynamics is viewed as ‘time evolution’ of the modes on ‘spatial hypersurfaces’. So one needs to look beyond null coordinates as they do not lead to a true Hamiltonian that can describe evolution of the modes. We note that ingoing field solutions $|\psi\rangle$ have a phase factor $e^{-i\omega u}$. Along a given ingoing null trajectory advanced null coordinate $v$ is constant. However, one can use retarded null coordinate $u$ to parameterize its propagation. In other words, ingoing field solutions $f_\omega(v)$, using the relation $v = u + 2r_s$, can be viewed as if $f_\omega(u) = e^{-i\omega u} f_\omega(0)$ where $u$ varies along the trajectory. Remarkably, this form can be compared with time evolution of a Schrodinger wave function $\psi_\omega(\tau) = e^{-i\omega \tau}\psi_\omega(0)$ for a mechanical system with energy $\omega$ and time coordinate $\tau$. We also know that a massless, free scalar field can be mapped into a set of harmonic oscillators by using Fourier transformation. These insights then suggest that we may define a timelike coordinate by slightly deforming retarded null coordinate $u$ and define a spacelike coordinate by deforming advanced null coordinate $v$ for an observer near past null infinity $\mathscr{I}^-$, say $\mathcal{O}^-$, as

$$\tau_\pm = t - (1 - \epsilon) r_s^\pm \quad \xi_\pm = -t - (1 + \epsilon) r_s, \tag{8}$$

where $\epsilon$ is a real-valued parameter. In general, one can choose the parameter in the domain $0 < \epsilon < 2$ such that $\tau_- \text{ and } \xi_-$ are timelike and spacelike coordinates respectively. Here we choose the parameter $\epsilon$ to be a small and positive such that $\epsilon > \epsilon^2$. This choice of parameter allows us to mimic the basic tenets of the Hawking effect very closely. In any case, final result will be independent of the explicit values of $\epsilon$. Similarly, we define another set of timelike and spacelike coordinates $\tau_+$ and $\xi_+$ as

$$\tau_+ = t + (1 - \epsilon) r_s^+ \quad \xi_+ = -t + (1 + \epsilon) r_s, \tag{9}$$

for an observer near $\mathscr{I}^+$, referred to as $\mathcal{O}^+$. We note that one can algebraically transform the two sets of the coordinates $|\mathcal{O}^\pm\rangle$ to each other by simply substituting $r_s \to -r_s$. We have suitably chosen the directions of $\xi_-$ and $\xi_+$ for later convenience (See FIG.1). Similar to the relation (5), one can derive an analogous relation $\xi_+ = \xi_- + 2r_s \ln(\xi_-/2r_s)$ which can be approximated in the domain $|\xi_-| \ll 2r_s$, as

$$\xi_- \approx 2r_s \epsilon^{\xi_-/2r_s}. \tag{10}$$

where $\xi_-$ and $\xi_+$ refer to the spatial coordinates on a $\tau_- = \text{constant}$ and $\tau_+ = \text{constant}$ surfaces for the observers $\mathcal{O}^-$ and $\mathcal{O}^+$ respectively (details of the canonical derivation is provided in an accompanying paper [15]).

Scalar Field Hamiltonian. The Hawking effect is crucially connected with the structure of the Schwarzschild
FIG. 1: (a) In Penrose diagram, shaded region represents the collapsing matter shell. Ingoing null rays depart from past null infinity $J^-$ whereas outgoing null rays arrive at future null infinity $J^+$. Near-null coordinates are $\tau_\pm$ and $\xi_\pm$. (c) Arrival time $\Delta \tau_k$ for $k^{th}$ mode in arbitrary units.

metric in the $t-r$ plane. So for simplicity now onward we consider only the 1+1 dimensional system. For both the observers $\mathcal{O}^-$ and $\mathcal{O}^+$, the metrics are of the form

$$ds^2 = g^\pm_{\mu\nu} dx^\mu dx^\nu = \frac{\Omega}{2} \left[-dt^2 \pm 2\epsilon d\tau \pm d\xi^2\right].$$ (11)

For large radial distances, the 4-dimensional scalar field action can be reduced to the form

$$S_\varphi = \int d\tau d\xi \left[-\frac{1}{2}\sqrt{-g} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi\right],$$

where $g^{\pm}_{\mu\nu}$ is flat but has off-diagonal terms. The corresponding scalar field Hamiltonians are

$$H^\pm_\varphi = \int d\tau d\xi \frac{1}{\epsilon} \left[\left(\frac{\Pi^2}{2} + \frac{1}{2}(\partial_\tau \varphi)^2\right) + \Pi \partial_\xi \varphi\right],$$ (12)

where the lapse function $N = 1/\epsilon$, the shift vector $N^1 = 1/\epsilon$ and the determinant of the spatial metric $g = 1$. The Poisson bracket between the field $\varphi$ and its conjugate momentum $\Pi$ can be written as

$$\{\varphi(\tau_\pm, \xi_\pm), \Pi(\tau_\pm, \xi'_\pm)\} = \delta(\xi_\pm - \xi'_\pm).$$ (13)

Using equations of motion, the field momentum can be expressed as $\Pi = \epsilon (\partial_\tau \varphi) - (\partial_\xi \varphi)$.

Fourier modes.- The Fourier modes of the scalar field can be defined for both observers as

$$\varphi = \frac{1}{\sqrt{V_\pm}} \sum_k \hat{\varphi}_k e^{ik\xi_\pm}; \quad \Pi = \frac{1}{\sqrt{V_\pm}} \sum_k \sqrt{q} \tilde{\pi}_k e^{ik\xi_\pm},$$ (14)

where $\hat{\varphi}_k = \hat{\varphi}_k(\tau_\pm), \tilde{\pi}_k = \tilde{\pi}_k(\tau_\pm)$ are the complex-valued mode functions. The spatial volume $V_\pm = \int d\xi_\pm \sqrt{q}$ are formally divergent. To avoid dealing with explicitly divergent quantity, we choose a fiducial box with finite volume. Then the wave-vectors are $k \in \{k_r\}$ where $r_\pm = 2\pi r/L_\pm$ with $r$ being a non-zero integer and $L_\pm$ being the length of the box. The scalar field Hamiltonian (12) can be expressed as $H^\pm_\varphi = \sum_k \frac{1}{2}(H^k_\varphi + D^k_\varphi)$ where Hamiltonian density for $k^{th}$ mode

$$H^k_\varphi = \frac{1}{2} \tilde{\pi}_k \tilde{\pi}_{-k} + \frac{1}{2}|k|^2 \hat{\varphi}_k \hat{\varphi}_{-k}$$

diffeomorphism generator

$$D^k_\varphi = -\frac{ik}{2} \left(\tilde{\pi}_k \hat{\varphi}_{-k} - \tilde{\pi}_{-k} \hat{\varphi}_k\right).$$

The associated Poisson bracket is $\{\hat{\varphi}_k, \tilde{\pi}_{-k}\} = \delta_k k$. We can relate the Fourier modes between the two different observers as

$$\hat{\varphi}_k = \sum_{k} \hat{\varphi}_k F_0(k, -\kappa) , \quad \tilde{\pi}_k = \sum_k \tilde{\pi}_k F_1(k, -\kappa)$$ (15)

where $\hat{\varphi}_k = \hat{\varphi}_k(\tau^0_k), \tilde{\pi}_k = \tilde{\pi}_k(\tau^0_k), \tilde{\pi}_k = \tilde{\pi}_k(\tau^0_k)$ and $\kappa = \kappa(\tau^0_k)$ [15]. The coefficient functions $F_m(k, \kappa)$ are similar to the Bogoliubov coefficients in covariant formulation and are likewise formally divergent. It is shown [11, 15] that these coefficients can be regularized to render them finite. The regulated coefficients $F^\delta_m(\pm|k|, \kappa)$ reduces to the exact expression when the regulator $\delta$ is removed i.e. $\lim_{\delta \to 0} F^\delta_m(k, \kappa) = F_m(k, \kappa)$ and satisfy following key relations [15]

$$F^\delta_0(-|k|, \kappa) = e^{2\pi \tau_\kappa - i \delta \tau} F^\delta_0(|k|, \kappa),$$

$$F^\delta_0(\pm|k|, \kappa) = \mp \frac{\kappa}{|k|} F^\delta_0(\pm|k|, \kappa).$$ (16)

Number operator.- In order to quantize the scalar field we follow the method as used in [16] where one canonically quantizes each Fourier mode. In particular, the expectation value of the Hamiltonian density operator of a positive frequency mode i.e. $\kappa > 0$, for the observer $\mathcal{O}^+$ in the vacuum state $|\_\rangle$ of the observer $\mathcal{O}^-$ i.e. $\langle \mathcal{H}^+_\varphi | \_\rangle = \langle \_ | \mathcal{H}^+_\varphi | \_ \rangle$ can be expressed as [15]

$$\langle \mathcal{H}^+_\varphi | \_ \rangle = e^{2\pi \kappa/\epsilon} + \frac{1}{\xi(1 + 2\delta)} \sum_{r=1}^\infty \frac{1}{r^{1+2\delta}} \langle \mathcal{H}^+_k \rangle,$$ (17)

where Riemann zeta function $\zeta(1 + 2\delta) = \sum_{r=1}^\infty r^{-(1+2\delta)}$ and we have used the properties of the vacuum state such that $\langle 0_k | \hat{\varphi}_k | 0_k \rangle = 0$ and $\langle 0_k | \tilde{\pi}_k | 0_k \rangle = 0$. We define the number density operator which represents the Hawking quanta as

$$\hat{N}_\kappa = \left[\mathcal{H}^+_\varphi - \lim_{\epsilon \to 0} \mathcal{H}^+_\varphi \right]|\kappa|^{-1},$$ (18)

which makes it amply clear that existence of these quanta are tied to the non-zero values of the surface gravity $\kappa$ at the horizon. Besides, this definition becomes crucial for the situation where the notion of creation and annihilation operators are not readily available like in polymer quantization.

The Fourier modes that we have considered so far, are in general complex valued functions. So in order to avoid double counting, as $\varphi$ is real-valued, here we make a choice by setting imaginary components of the modes $\hat{\phi}^*_k = 0$ and $\pi^*_k = 0$. This leads diffeomorphism generator...
$D_k$ to vanish identically. Further, by redefining the real part of the modes as $\phi_k \equiv \phi_k^e$ and $\pi_k \equiv \pi_k^e$, we can reduce the Hamiltonian density to its regular harmonic oscillator form $\mathcal{H}_k = \frac{1}{2} \pi_k^2 + \frac{1}{2} k^2 \phi_k^2$ along with the Poisson bracket $\{\phi_k, \pi_k\} = \delta_{k,k'}$. The corresponding energy spectrum can be written as $\mathcal{H}_k|_{n_k} = (n + \frac{1}{2})|k||n_k|$ where $n \geq 0$. Therefore, in Fock quantization, $\mathcal{H}_k = \frac{1}{2}|k|$ for all modes. The Eqn. (17) and (18) together then imply

$$N_\omega \equiv \langle \tilde{N}_{k=\omega} \rangle = \frac{1}{e^{2\pi \omega/\kappa} - 1} = \frac{1}{e^{(4\pi r_s)\omega} - 1}, \quad (19)$$

which corresponds to a thermal spectrum at Hawking temperature $T_H = \kappa/(2\pi k_B) = 1/(4\pi r_s k_B)$. It shows that we can derive the exact thermal spectrum of the Hawking effect also using Hamiltonian formulation.

**Polymer quantization.**– In polymer quantization, energy eigenvalues for the $k$th oscillator is given by [10]

$$E_k^{2n} = \frac{1}{4g} + \frac{g}{2} A_n(g), \quad E_k^{2n+1} = \frac{1}{4g} + \frac{g}{2} B_{n+1}(g), \quad (20)$$

where $n \geq 0$, $A_n$ and $B_n$ are Mathieu characteristic value functions. The dimensionless parameter $g \equiv |k| l_*$ where $l_*$ is the polymer length scale. For small $g$, the energy spectrum (20) reduces to $E_k^{2n} = 1/4g + O(g^3)$. It implies that polymer quantization correctly reproduces the spectrum for sub-Planckian modes. However, significant non-perturbative modifications in the spectrum are seen for super-Planckian modes. In particular, for large $g$, ground state energy can be approximated as $E_0/k = 1/4g + O(g^{-3})$. So unlike in Fock quantization where $\langle \mathcal{H}_k \rangle/k = \frac{1}{2}$ for all $k$, in polymer quantization $\langle \mathcal{H}_k \rangle/k \to 0$ for the trans-Planckian modes as $k \to \infty$. Therefore, when one removes the regulator $\delta$ in polymer quantization, the expectation value of the number operator $\langle \hat{N} \rangle_{poly}$, due to the form of Eqn. (17) and the zeta function identity $\lim_{s \to 0} \zeta(1 + s) = 1$, becomes

$$N_{poly} = \langle \hat{N}_{poly} \rangle = 0. \quad (21)$$

This property of the number operator, having same mathematical expression, is identical to the case of Unruh effect [11]. Therefore, an asymptotic future observer $\mathcal{O}^+$ would not perceive any Hawking quanta in polymer quantization, in contrary to the Fock quantization.

**Duration of Hawking effect.**– We note that expectation value of number operator is discontinuous in the limit $l_* \to 0$. In order to physically understand this behavior, let us consider a future observer located at a fixed $r \gg r_s$. The proper time interval $\Delta \tau$ for this observer then follows the relation $\Delta \tau = \Delta t = \Delta t_+$, as fixed $r$ implies $\Delta r_+ = -\Delta \xi$. For this observer, the difference in arrival time for two modes which were emitted from the coordinate points $\xi^1$ and $\xi^2$ on a fixed $\tau_-$ surface near $\mathcal{I}^-$ can be written using the relation (10) as $\Delta t = 2r_s \ln (\xi^1/\xi^2)$. The relevant modes for Hawking radiation are emitted from the region $|\xi_-| < 2r_s$. So for simplicity we choose the arrival time of the mode emitted from $\xi_- = 2r_s$ as the ‘beginning’ of Hawking radiation. Then the difference in arrival time for a mode emitted from a general coordinate point $\xi_- (\leq 2r_s)$ would be

$$\Delta \tau = 2r_s \ln (2r_s / \xi^-). \quad (22)$$

Clearly, the modes whose point of emission $\xi^-$ is closer to the origin, arrive later near $\mathcal{I}^+$. For a relativistic mode in ground state, one can always associate a de-Broglie wavelength (like ‘width’ of the quanta) as $\lambda_k = h/(\kappa E_k^0)$. Naturally, point of emission cannot be made more accurate than its de-Broglie wavelength. Hence we choose closest possible point of emission for the $k$th mode as $\xi^- = \lambda_k^0$. It may be checked that the proper wavelength of these modes at the time of arrival near $\mathcal{I}^+$ would be $\lambda_k^0 \sim 2r_s$ which can be viewed as the ‘Wien’s displacement law’ for Hawking radiation. We now define the ‘duration’ of the Hawking effect to be the arrival time of the mode with least possible de-Broglie wavelength, given by

$$\tau^H \equiv \max (\Delta \tau_e) = \lim_{k \to \infty} 2r_s \ln \left[ \frac{r_s}{\pi} \langle \mathcal{H}_k \rangle \right]. \quad (23)$$

In Fock quantization $\langle \mathcal{H}_k \rangle = \frac{1}{2}|k|$ for all modes, including trans-Planckian modes, which implies $\tau^H \to \infty$. On the contrary, for trans-Planckian modes in polymer quantization $\langle \mathcal{H}_k \rangle \approx \frac{1}{\kappa}$ which implies $\tau^H \approx 2r_s \ln (r_s/4\pi l_*)$ (see FIG 1). If we take $l_*$ to be Planck length then for a solar mass black hole this duration $\tau^H \approx 1.7$ milliseconds whereas for an ultra-massive black hole with mass $M = 4 \times 10^{10} M_\odot$ (like one at the center of the galaxy S5 0014+81) the duration $\tau^H \approx 2.8$ years. This short duration of the Hawking effect explains why an observer in asymptotic future would not perceive any Hawking quanta in polymer quantization.

**Information loss paradox.**– In Fock quantization, the Hawking effect persists ad infinitum. Therefore, one argues that it would eventually lead to a complete evaporation of the black hole. This in turn leads to the so-called information loss paradox [11,5], as from thermal radiation alone one cannot recover information about the collapsing matter shell which led to the formation of the black hole. Recently Unruh and Wald have classified the proposals to resolve information loss paradox in four categories: (I) fuzzball formation (II) firewall scenario (III) Planckian remnant and (IV) Planckian final burst [3]. However, in the scenario as implied by the polymer quantization, the Hawking radiation stops after a short duration leaving the classical black hole unchanged. Consequently, there is no loss of information. Therefore, it provides a new way to resolve the information loss paradox and it requires modification only in the trans-Planckian physics.
Discussions.-- In summary, we have shown that the Hawking effect is short-lived in polymer quantization of matter field and it eventually disappears to an asymptotic future observer. In order to arrive at these results we have introduced a set of near-null coordinates which allowed us to have an exact derivation of the Hawking effect using Hamiltonian formulation. In polymer quantization, the duration of the Hawking effect would appear to be few milliseconds for a solar mass black hole whereas it would be few years for an ultra-massive black hole. These predictions are testable in principle and may allow one to verify or rule out the given hypothesis. Furthermore, this short-lived Hawking effect scenario provides a new way to resolve the so-called information loss paradox.

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