

Hawking Evaporation and the Fate of Black Holes in Loop Quantum Gravity

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A recent covariant formulation, that includes nonperturbative effects from loop quantum gravity (LQG) as self-consistent effective models, has revealed the possibility of nonsingular black hole solutions. The new framework makes it possible to couple scalar matter to such LQG black holes and derive Hawking radiation in the presence of quantum spacetime effects while respecting general covariance. Standard methods to derive particle production both within the geometric optics approximation and the Parikh–Wilczek tunnelling approach are therefore available and confirm the thermal nature of Hawking radiation. The covariant description of scale-dependent decreasing holonomy corrections maintains Hawking temperature as well as universality of the low-energy transmission coefficients, stating that the absorption rates are proportional to the horizon area at leading order. Quantum-geometry effects enter the thermal distribution only through subleading corrections in the graybody factors. Nevertheless, they do impact energy emission of the black hole and its final state in a crucial way regarding one of the main questions of black-hole evaporation: whether a black-to-white-hole transition, or a stable remnant, is preferred. For the first time, a first-principles derivation, based on a discussion of backreaction, finds evidence that points to the former outcome.

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Introduction—Hawking’s celebrated black-hole information loss paradox has been instrumental in fueling research in fundamental physics over the last half a century. A black hole, which forms when a sufficiently massive star collapses onto itself, has an event horizon that effectively isolates it from the rest of the universe. However, it was soon understood that, quantum-mechanically, black holes behave as thermal objects, emitting blackbody radiation as viewed by an external observer. This observation immediately gives rise to a paradox. What happened to the pure state of the star whose collapse formed the black hole if, in the end, it completely evaporates by Hawking radiation? Two effects, ignored in Hawking’s calculation, can help resolve this question. First, backreaction of the radiation stress-energy tensor on the spacetime geometry, although small for

large-mass black holes, could become non-negligible at later stages of the evaporation process as the mass shrinks to zero. Second, at the final stage, backreaction alone is insufficient, and quantum-gravity corrections must be considered. Since Hawking evaporation in classical spacetime ends at a singularity, complicating any discussion of backreaction effects on the shrinking black hole, potential regularization effects from quantum gravity, provided they are sufficiently well understood and preserve a meaningful and covariant spacetime interpretation, may make it possible to address these questions. This is our primary motivation.

Vacuum spherically symmetric solutions of loop quantum gravity (LQG) have hinted at resolving the singularity at the core of black holes by a quantum “transition surface” [1–3]. Nevertheless, self-consistent dynamical solutions must be examined to understand the fate of stellar collapse and, thus, how the information-loss paradox is resolved in LQG. In this Letter, we study Hawking radiation from a covariant LQG black-hole model with scale-dependent decreasing holonomy functions, a refined version of what is commonly referred to as the $\bar{\mu}$ -scheme [4]. While we recover a thermal distribution with the same asymptotic temperature as in general relativity, $T_H \sim 1/M$, LQG parameters related to spacetime discreteness affect the graybody factors and, hence, the emission rate. This new effect provides a mechanism for thermal stabilization and suggests, when

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taken in isolation, the formation of a remnant at a sub-Planckian mass. However, horizon fluctuations as well as new features of the quantum gravitational force at high density, that destabilize the system gravitationally, must also be taken into account. These crucial ingredients are derived from a covariant spacetime description of LQG modifications, combined with a qualitative discussion of backreaction in the presence of matter couplings, as well as implications of the uncertainty relation. The derived mass hierarchy of the three central effects—thermal stabilization, significant horizon fluctuations, and gravitational destabilization—shows that a black-to-white-hole transition is preferred in LQG over the formation of a remnant. Our conclusion requires a careful consideration of consistency conditions that have only recently become available for spacetime models including quantum-geometry corrections. To simplify the discussion, we use Planck units where $c = \hbar = k_B = G = 1$.

Effective LQG—Previous attempts to develop effective spacetimes that incorporate holonomy corrections have failed to satisfy general covariance. Some approaches use loop quantum cosmology techniques to model the interior region of black holes and then extrapolate the results to the exterior region [5,6]. Others involve partial gauge fixing of the system prior to modifying the dynamics [7,8]. While these approaches arguably succeed in resolving the singularity issue, they lack the gauge freedom necessary to test their covariance status, and thus it is difficult to extract robust, meaningful physical effects from them.

Emergent modified gravity [9,10] is a Hamiltonian approach to gravity that incorporates spacetime not as a presupposed quantity but rather as an object derived (or emerging) from a set of covariance conditions. This key idea results in weaker assumptions on the basic setup of a gravitational system and evades previous no-go theorems [11] about possible covariant realizations of central LQG effects. As a result, new theories are now available which, in particular, include modification functions in spherically symmetric models that play the role of holonomy corrections of LQG [3,10]. These theories use standard phase-space variables of canonical spherically symmetric gravity and impose a modified constraint that is required to preserve anomaly-free Poisson brackets and to generate gauge transformations corresponding directly to spacetime diffeomorphisms when equations of motion are used (on shell). These conditions imply a set of differential equations to be satisfied by the constraint which has been fully solved in Ref. [10] up to the classical order in spatial derivatives. An analysis of canonical transformations reveals that any modified constraint can be brought to a form suitable for loop quantization; see the End Matter for more details. Moreover, the canonical formulation facilitates a combination of emergent modified gravity with an expansion of the metric by quantum fluctuations along the lines of Ref. [12]. While details of this extension will not be necessary for our

qualitative discussion, it shows that quantum fluctuations remain under control in this framework.

In the Schwarzschild gauge, solving these effective LQG equations of motion results in the line element

$$ds^2 = -f(x)dt^2 + \frac{dx^2}{h(x)f(x)} + x^2 d\Omega^2, \quad (1)$$

where $f(x) = 1 - 2M/x$, $h(x) = 1 + \lambda(x)^2 f(x)$, and $\lambda(x)$ is the holonomy function with classical limit $\lambda \rightarrow 0$. Here, we restrict discussions to a decreasing function $\lambda(x) = \sqrt{\Delta}/x$, where $\Delta \propto \ell_P^2$ may be interpreted as a discreteness parameter in LQG, mimicking the traditional $\bar{\mu}$ -scheme [4]. [The analysis for a general $\lambda(x)$ can be found in Ref. [13].] The line element (1) preserves a horizon at the classical location $x = 2M =: x_H$. It contains an additional coordinate singularity at $x = x_\Delta < x_H$ defined as the solution to $h(x_\Delta) = 0$ and given by

$$x_\Delta = 3M\delta^{1/3} \frac{(1 + \sqrt{1 + \delta})^{2/3} - \delta^{1/3}}{(1 + \sqrt{1 + \delta})^{1/3}}, \quad (2)$$

where $\delta = \Delta/(27M^2)$, such that $x_\Delta(M) \approx (2M\Delta)^{1/3}$ for $M \gg \sqrt{\Delta}$. This expression can be thought of as the minimal areal radius for the vacuum solution which marks a reflection-symmetry surface [14] joining the black hole to a white hole; see Ref. [3] for details. (Qualitatively, these properties are similar to the case of constant λ given earlier in [1,2]; see also Ref. [15].)

A massless test scalar field can be minimally coupled to the holonomy-modified system such that its propagation follows the standard Klein–Gordon equation [16–18]

$$\tilde{g}^{\mu\nu} \tilde{\nabla}_\mu \tilde{\nabla}_\nu \phi = 0, \quad (3)$$

where the derivative operator $\tilde{\nabla}$ is compatible with the modified background metric $\tilde{g}_{\mu\nu}$ defined by Eq. (1). However, there are other couplings allowed by covariance, which may be relevant for effective descriptions of new quantum-gravity effects. We show an example at the end of this Letter. The importance of imposing strict covariance conditions in deriving the line element (1) in LQG, and while coupling scalar matter to it, will also be seen in what follows. These are key new ingredients of our construction.

Thermal distribution—Our first derivation determines tunnelling rates for Hawking particles in the Parikh–Wilczek picture [19]. The metric is important because it may modify the potential and corresponding tunnelling rates. We present details for a derivation with our emergent line element in Supplemental Material (SM) [20], but note one crucial feature here: in the metric coefficients, λ enters only through the function $h(x)$, where it is multiplied by

$f(x)$. The latter function vanishes at the horizon, right where the tunnelling process is relevant. As a consequence, the standard tunnelling rate is unmodified in this specific version of an effective line element. [Mathematically, strict equality of the tunnelling rates on classical and quantum backgrounds is implied by the application of residues in an action integration with poles at the horizon. The residue does not depend on λ because this function disappears when the expressions are evaluated at the horizon, where $f(x) = 0$.] With an unmodified tunnelling rate, the standard expressions for Hawking temperature and Bekenstein–Hawking entropy hold true as well, even in the presence of holonomy corrections. This outcome is in contrast to incomplete proposals of LQG effects in black-hole geometries which had not imposed all the necessary consistency conditions [21–25]. It corrects a long-standing misconception [26] by showing that Hawking entropy, as measured by an asymptotic observer, remains the same in LQG. We now turn to the question of which aspects of Hawking evaporation are affected by LQG effects.

Hawking [27] was the first to analyze the behavior of a scalar field in the vicinity of a black hole while neglecting backreaction. The distant past, before the cloud of dust had collapsed, and the distant future, after the black hole has formed as seen by an asymptotic observer, are both approximated as flat space. Matching incoming with outgoing modes through the black-hole geometry provides the spectrum of Hawking radiation. Since the curved geometry is position dependent, this process happens in an effective potential which imposes a barrier that is easier to overcome for high-frequency modes. The Hawking spectrum of the frequency-dependent number of particles,

$$\langle N_\omega \rangle = \frac{\mathcal{T}_l(\omega)}{\exp(\omega T_H) - 1}, \quad (4)$$

therefore contains not only the temperature-dependent blackbody factor but also a graybody factor $\mathcal{T}_l(\omega)$. Unlike the former, the latter contains information about the emergent geometry because it is sensitive to a region around the horizon.

A test scalar field propagating in a spherically symmetric spacetime background described by (1) can be decomposed into radial and angular parts as $\phi(x^\mu) = \sum_{lm} e^{i\omega t} \Phi_{lm}(x) Y_{lm}(\theta, \varphi)$, with spherical harmonics $Y_{lm}(\theta, \varphi)$. Equation (3) then implies that the radial modes $\Phi(x)$ satisfy the equation

$$\left[\frac{d^2}{dx_*^2} + U_l(x) \right] (x \Phi_{lm}) = 0, \quad (5)$$

where x_* is the radial tortoise coordinate defined by $dx_* = f(x)^{-1} h(x)^{-1/2} dx$ and

$$\begin{aligned} U_l(x) &= \omega^2 - V_l(x) \\ &= \omega^2 - \left(1 - \frac{2M}{x} \right) \left[\frac{2M}{x^3} + \frac{l(l+1)}{x^2} \right. \\ &\quad \left. + \left(1 - \frac{2M}{x} \right) \left(\frac{3M\lambda(x)^2}{x^3} + \frac{\lambda(x)\lambda'(x)}{x} \left\{ 1 - \frac{2M}{x} \right\} \right) \right] \end{aligned} \quad (6)$$

is the effective potential. Equation (5) has a complicated solution, but it can still provide analytical information about physical effects in the asymptotic region. As before, compared with the classical expression, all λ -terms are multiplied by additional factors of $1 - 2M/x$ that vanish at the horizon. Nevertheless, graybody factors depend on λ because they are sensitive to the entire potential.

Analytical computations of graybody factors are not feasible [28,29], which complicates the precise evaluation of the evaporation process. We resort to the standard approximation of low-frequency modes in which we assume that the incoming scalar particle's energy is below the background thermal energy, $\omega \ll T_H$ or $M\omega \ll 1$, and expand the distribution in terms of the latter quantity. This step allows us to extract physical consequences.

In order to apply this approximation [30,31], we need to identify three regions: (i) close to the horizon where the potential is nearly vanishing, $V_l(x) \ll \omega^2$; (ii) the intermediate region near the peak of the potential where $V_l(x) \gg \omega^2$; and (iii) the asymptotic region where the potential flattens out again. By solving Eq. (5) in these three regions and requiring continuity at the boundaries, one obtains the wave function parametrized by the coefficient of the first region. The peak of the potential increases with the multipole l . Accordingly, high- l modes with low energy are mostly reflected back by the potential and hence can be neglected in the transmission coefficient. The resulting transmission coefficient is dominated by s -wave modes [28]. (For the same reason, it is sufficient to discuss scalar particles because higher spins do not contribute to the $l = 0$ mode.)

The part of the wave function that traverses the potential and enters the black hole can be quantitatively computed by evaluating the transmission coefficient

$$\mathcal{T}_0(\omega) = \frac{J_H}{J_{\text{in}}} = \frac{16M^2\omega^2}{(1 - 4M^2\omega^2)^2 + 4M^2\omega^2 \left(1 + \frac{\Delta}{12M^2} \right)^2},$$

where J_{in} and J_H denote the complete s -wave flux currents from the incident and transmitted waves, respectively. We recover universality of the black-hole absorption rate [32,33], which to leading order in $M\omega$ depends solely on the horizon area (without the need to impose the classical limit $\Delta \rightarrow 0$).

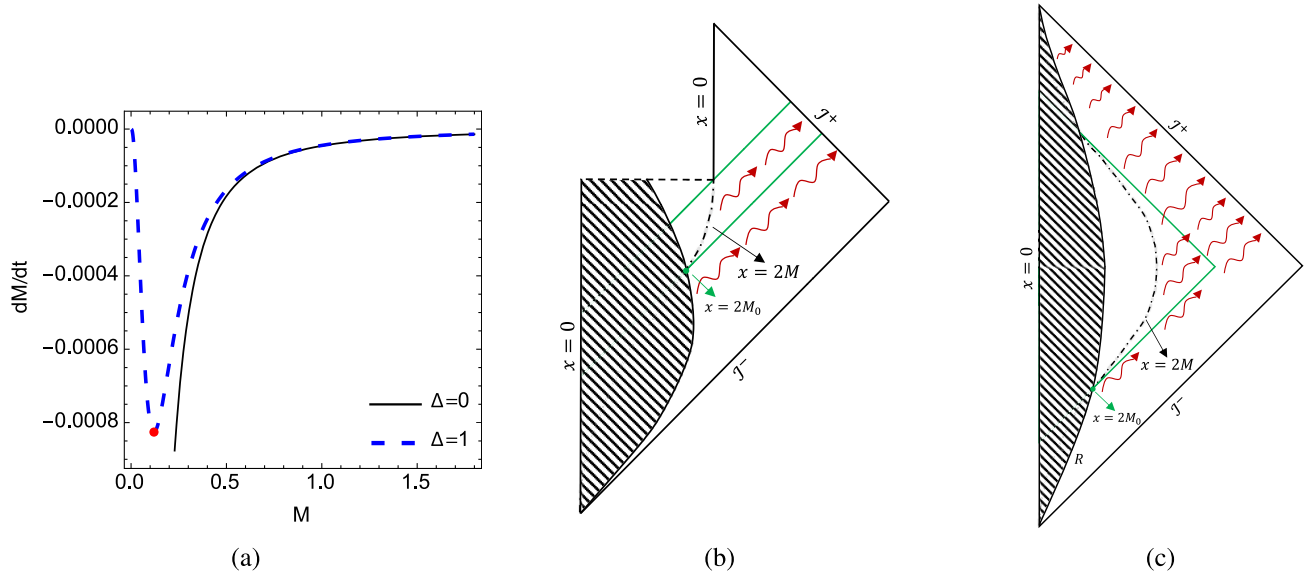


FIG. 1. (a) Energy emission for a classical black hole (solid line) and the holonomy corrected one (dashed line) near Planckian scales. (b) Classical collapse and evaporation. The black hole features a singularity denoted by the dashed line and evaporation ends in an empty Minkowski space. (c) Black-to-white-hole transition. The apparent horizon $x = 2M(t)$ is timelike due to energy loss; in the early stages it remains near null, but as energy emissions grow it deviates prominently into a timelike trajectory from which matter can escape if not met with a singularity.

Emission rate—Holonomy corrections modify the transmission amplitudes and, consequently, impact the energy emission rate:

$$\frac{dM}{dt} = -\frac{1}{2\pi} \int_0^\infty d\omega \frac{\omega T_0(\omega)}{\exp(8\pi M\omega) - 1}. \quad (7)$$

Backreaction is incorporated by a dynamical correction of the mass in the line element to account for the energy carried away by the outgoing modes. A detailed discussion in terms of energy-conserving particle tunnelling [19] can be found in Supplemental Material [20].

If general relativity was valid throughout the entire process, the black hole’s energy would eventually be fully depleted, returning to flat Minkowski spacetime [Fig. 1(b)]. Holonomy corrections from nonzero Δ enter as a subleading term that slows down the emission rate as the hole evaporates, suggesting that quantum-gravity corrections alter the final stage of a black hole compared with the classical theory. We arrive at this conclusion without setting a Planckian cutoff by hand, either in the momentum integral [34,35] (which would alter the emission rate by violating Lorentz invariance) or in the two-point function [36]. Instead, our result follows naturally from a covariant description of nonperturbative effects in LQG.

As shown in Fig. 1(a), quantum spacetime effects alter the energy emission rates of the black hole. While a classical black hole evaporates perpetually with an accelerating rate of energy emission, the new holonomy effect eventually slows down the evaporation process, around the

sub-Planckian mass $M_r \approx 0.12\sqrt{\Delta}$. After reaching this mass, the emission rate decreases sharply and its Hawking radiation slows down remarkably as compared to that of a classical black hole, achieving a vanishing mass only in an infinite amount of time. Quantum effects can therefore stabilize a black hole thermally, which, taken at face-value, seems to indicate the formation of a remnant.

Qualitatively, this outcome is not unexpected. Vacuum spherically symmetric solutions of effective LQG equations have been used to show that, near the Planck scale, the classical singularity is replaced by a reflection-symmetry surface. The eternal solution (1) suggests that once matter passes through this surface, it “bounces” into a white hole structure [1–3]. However, this vacuum description does not determine how the process of collapsing matter is perceived by a distant observer. Matter could bounce into a white hole which shares the same asymptotic region as the black hole and thus be visible to an observer or split off as a remnant with a disconnected asymptotic region. The latter case would be described as the formation of an interuniversal wormhole [3,37], and the former as a black-to-white-hole transition [38,39]. As we show in what follows, backreaction of Hawking radiation plays a central role in resolving this ambiguity of outcomes. We discuss the two potential final stages of the black hole: one described by a thermally stabilized remnant and the other by a transition process from a black hole to a white hole; see Fig. 1.

The evaporation process—When the black hole mass is large, $M \gg \sqrt{\Delta}$, the transmission coefficient (7) is

dominated by the classical terms, and the leading term in (7) implies an evaporation time going as M^3 . The prediction remains robust near Planckian mass until it deviates from the M^3 behavior near M_r ; see Fig. 1(a).

The emission rate slows down at a sub-Planckian mass M_r , suggesting a stable remnant. However, once in this quantum spacetime regime, the horizon area can only be resolved up to quantum fluctuations. Hawking radiation then ends when the horizon area reaches, within quantum fluctuations, the smallest possible radius $x_\Delta(M)$. Based on the uncertainty relation, quantum fluctuations should be proportional to $\sqrt{\hbar} \propto \ell_{\text{Pl}}$, with an additional factor of M required for the correct dimensions of area fluctuations. Following Ref. [40], a specific expression for area fluctuations is given by $2M\sqrt{8\pi\Delta}$, where we have replaced ℓ_{Pl} by $\sqrt{\Delta}$ as the fundamental UV-scale in LQG. Upon equating the fluctuating horizon radius with the size of the minimum-area surface, $x_\Delta(M)$, we find $M_f \approx 4.16\sqrt{\Delta}$, a value much bigger than the mass M_r when the emission rate starts decreasing to initiate the formation of the remnant. The dissolving horizon implies an end of Hawking evaporation at M_f rather than M_r .

Backreaction can be taken into account by adjusting the black-hole geometry to decreasing masses according to (7). When the evaporation process ends at M_f , the continuum spacetime geometry, and therefore the classical picture of a horizon-bound black hole, dissolves. Instead, we have a superposition of a horizonless object and the final stage of the collapsing black hole. Semiclassically, this quantum state and its future gives a coarse-grained description of a transition from a black hole to its opened-up interior. Importantly, this stage is approached in finite time, before a remnant could form.

In this process, energy of the collapsing matter leaks out after some time as described in Fig. 1(c). A final observation made in Ref. [18] allows us to distinguish between a wormhole outcome and a white hole: some quasinormal modes of a nonminimally coupled scalar field—derived from the same background (1)—can become unstable for Planckian black holes at the critical mass $M_c \approx 0.57\sqrt{\Delta} > M_r$, jump-starting the white-hole emission. We obtain the same thermal stabilization for Planckian black holes for a nonminimally coupled scalar field, as shown in Supplemental Material [20], but it destabilizes gravitationally. This outcome can be understood as an implication of a repulsive quantum-gravity contribution, which helps to avoid the singularity and can be relevant for Planckian black holes in their exterior, too. A quantum black hole is no longer gravitationally bound and therefore transitions to a white hole. This instability does not happen for minimal coupling, but insisting on such a choice within a general effective theory of quantum gravity would amount to fine-tuning. Even if nonminimal coupling coefficients are small, the instability of some

modes ensures that their implications are relevant. Instead of thermal Hawking radiation, the final stages of black-hole evaporation are dominated by a gravitational instability governed by unitary quantum evolution.

Conclusions—As the mass of an evaporating LQG black hole shrinks toward M_r , holonomy corrections dramatically slow down Hawking radiation. When the horizon area reaches a Planckian size, quantum fluctuations can no longer be ignored. Within fluctuations, the horizon reaches the minimum-area reflection-symmetry surface at $M_f > M_r$. In the underlying quantum state, there is a finite probability for a sub-Planckian black hole to tunnel into a white hole [41], an outcome that would always be preferred given that it takes a long time to form the remnant at mass M_r . Going beyond heuristic claims, we provide a physical mechanism, based on a holonomy-induced instability in a quasinormal mode which constitutes an explicit candidate for an initial white-hole emission at a mass scale $M_c > M_r$. Crucially, the mass hierarchy $M_r < M_c < M_f$ derived here reveals that the black hole first turns quantum, as shown by significant quantum fluctuations at M_f , and forms a superposition of a horizon-covered black hole with an exposed white hole. This stage is followed by the quantum-gravitational instability triggering an explosion at M_c , well before Hawking radiation starts weakening at M_r . Although the precise values of these masses may change if quantum fluctuations are included in the metric [12], our scenario would remain valid since it relies only on generic quantum features. Even if the inner geometry changes drastically due to quantum fluctuations and M_r disappears in a fuzzy nonclassical background, our final conclusion of not having a remnant remains robust.

Our results provide the first covariant, self-consistent description of black-hole evaporation in LQG with several crucial ingredients of quantum spacetime falling into the right places. Compared to previous attempts, we incorporated two new contributions with important implications for the physical spacetime of LQG black holes: consistent matter couplings and backreaction effects. Computing Hawking radiation for nondynamical vacuum models initially suggests the formation of a stable remnant. However, upon factoring in backreaction and fluctuation effects, we find that the black hole transitions to a white hole before the remnant mass is reached, enabled by new (quantum) gravitational instabilities at high density. Thus, dynamical evaporation is essential not only for differentiating between competing scenarios but also in establishing the final stage of black holes in quantum gravity. The inclusion of backreaction is possible due to the singularity-free nature of our dynamical solutions. In this picture, information inside the black hole naturally flows into the white hole, but only once the mass of the evaporating black hole is comparable to Planckian scales where LQG effects become large and backreaction cannot be ignored.

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Data availability—No data were created or analyzed in this Letter.

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End Matter

Emergent modified gravity—In the classical spherically symmetric theory, the gravitational momenta E^x and E^φ are components of the densitized triad directly related to the ADM line element,

$$ds^2 = -N^2 dt^2 + q_{xx}(dx + N^x dt) + q_{\varphi\varphi} d\Omega^2, \quad (\text{A1})$$

by $q_{xx} = (E^\varphi)^2/E^x$ and $q_{\varphi\varphi} = E^x$, while the configuration variables are directly related to the extrinsic-curvature components $\mathcal{K}_\varphi = K_\varphi$ and $\mathcal{K}_x = 2K_x$. In emergent modified gravity, a modified Hamiltonian constraint $\tilde{H}$ is considered while the phase-space structure of the classical theory is preserved as well as the diffeomorphism constraint

$$H_x = E^\varphi K'_\varphi - K_x (E^x)'. \quad (\text{A2})$$

However, instead of presupposing the relation between the phase-space variables and spacetime, it is derived from the requirement that the modified constraint satisfies the hypersurface-deformation brackets

$$\{H_x[N^x], H_x[\epsilon^x]\} = -H_x[\mathcal{L}_{\epsilon^x} N^x], \quad (\text{A3})$$

$$\{\tilde{H}[N], H_x[\epsilon^x]\} = -\tilde{H}[\epsilon^x N'], \quad (\text{A4})$$

$$\{\tilde{H}[N], \tilde{H}[\epsilon^0]\} = -H_x[\tilde{q}^{xx}(\epsilon^0 N' - N\{\epsilon^0\}')]. \quad (\text{A5})$$

with the structure function $\tilde{q}^{xx}$ possibly different from its classical expression $q^{xx} = q_{xx}^{-1} = E^x/(E^\varphi)^2$. The inverse of this modified structure function, $\tilde{q}_{xx} = 1/\tilde{q}^{xx}$, now plays the role of the radial component of the line element (A1). A final condition is imposed for the gauge transformation of the modified spacetime metric $\tilde{g}_{\mu\nu}$, as generated by the modified constraint, to correspond to infinitesimal coordinate transformations (on shell),

$$\delta_\epsilon \tilde{g}_{\mu\nu} = \mathcal{L}_\epsilon \tilde{g}_{\mu\nu}. \quad (\text{A6})$$

According to emergent modified gravity [9,10], the most general, covariant Hamiltonian constraint in a spherically symmetric vacuum, up to second-order derivatives and quadratic in first-order derivative terms, can be derived from solving the covariance conditions. Here, we will consider the special class of constraints given by [10]

$$\begin{aligned} \tilde{H}_{\text{grav}} = & -\frac{\sqrt{E^x}}{2} \left[E^\varphi \left(\frac{1}{E^x} + \left\{ \frac{1}{E^x} - 2 \frac{\partial \ln \lambda^2}{\partial E^x} \right\} \frac{\sin^2(\lambda K_\varphi)}{\lambda^2} + 4 \left\{ \frac{K_x}{E^\varphi} + \frac{K_\varphi}{2} \frac{\partial \ln \lambda^2}{\partial E^x} \right\} \frac{\sin(2\lambda K_\varphi)}{2\lambda} \right) \right. \\ & \left. + \frac{(\{E^x\}')^2}{E^\varphi} \left(\left\{ \frac{K_x}{E^\varphi} + \frac{K_\varphi}{2} \frac{\partial \ln \lambda^2}{\partial E^x} \right\} \lambda^2 \frac{\sin(2\lambda K_\varphi)}{2\lambda} - \frac{\cos^2(\lambda K_\varphi)}{4E^x} \right) + \left(\frac{(E^x)'(E^\varphi)'}{(E^\varphi)^2} - \frac{(E^x)''}{E^\varphi} \right) \cos^2(\lambda K_\varphi) \right], \end{aligned} \quad (\text{A7})$$

with the associated structure function

$$\begin{aligned} \tilde{q}^{xx} &= \left(1 + \lambda^2 \left(\frac{(E^x)'}{2E^\varphi} \right)^2 \right) \cos^2(\lambda K_\varphi) \frac{E^x}{(E^\varphi)^2} \\ &=: h \frac{E^x}{(E^\varphi)^2}, \end{aligned} \quad (\text{A8})$$

where $\lambda = \sqrt{\Delta/E^x}$ with positive constant Δ . The classical constraint is recovered in the limit $\Delta \rightarrow 0$. The emergent line element is given by

$$ds^2 = -N^2 dt^2 + \tilde{q}_{xx}(dx + N^x dt)^2 + E^x d\Omega^2, \quad (\text{A9})$$

where $\tilde{q}_{xx} = 1/\tilde{q}^{xx}$. While (A7) is a special case, its dependence on the curvature variable K_φ is unique (up to

canonical transformations) according to the general expression derived in Ref. [10].

Scalar matter can be coupled to emergent modified gravity by adding $P_\phi \phi'$ to the diffeomorphism constraint (A2), and a Hamiltonian constraint contribution $\tilde{H}_\phi$ to (A7) such that the algebra (A3)–(A5) and the spacetime covariance condition (A10) are preserved. Furthermore, the matter covariance condition

$$\delta_\epsilon \phi = \mathcal{L}_\epsilon \phi \quad (\text{A10})$$

(on shell) is imposed, too. The minimal coupling to emergent modified gravity is realized by the contribution

$$\tilde{H}_\phi^{(\text{MC})} = \frac{\sqrt{\tilde{q}^{xx}}}{2} \left[\frac{P_\phi^2}{E^x} + E^x (\phi')^2 \right] + \frac{E^x}{2} \sqrt{\tilde{q}_{xx}} V, \quad (\text{A11})$$

and the nonminimal coupling is realized by

$$\begin{aligned} \tilde{H}_\phi^{(\text{NMC})} &= \frac{P_\phi^2}{E^\varphi} \frac{1}{2\sqrt{E^x}} \left(1 + \lambda^2 \left(\frac{(E^x)'}{2E^\varphi} \right)^2 \right) \cos^2(\lambda K_\varphi) \\ &+ \frac{(E^x)^{3/2}}{2E^\varphi} (\phi')^2 + \sqrt{E^x} E^\varphi V. \end{aligned} \quad (\text{A12})$$

Angular momentum contributions to the constraint from a nonspherical test scalar field are incorporated as part of the potential V as a result of the standard spherical-harmonics decomposition [18].

The canonical transformation

$$\begin{aligned} K_\varphi &= \frac{\bar{\lambda}}{\lambda} \bar{K}_\varphi, & E^\varphi &= \frac{\bar{\lambda}}{\lambda} \bar{E}^\varphi \\ K_x &= \bar{K}_x - \bar{E}^\varphi \bar{K}_\varphi \frac{\partial \ln \lambda}{\partial E^x}, \end{aligned} \quad (\text{A13})$$

where the new phase-space variables are written with a bar (the variables E^x , ϕ , and P_ϕ remain unchanged in this transformation) and constant $\bar{\lambda}$, turns the modified Hamiltonian constraint fully periodic in $\bar{K}_\varphi$ and hence suitable for loop quantization by standard means. The uniqueness of this dependence singles out this constraint as the only covariant effective loop quantum gravity model for black holes. A generalization of the constraint further allows for point-holonomy modifications of the scalar field [17].

Toward quantum theory—Upon the canonical transformation (A13) the Hamiltonian can be written as

$$\begin{aligned} \tilde{H}_{\text{grav}} &= -\frac{\bar{\lambda}}{\lambda} \frac{\sqrt{E^x}}{2} \left[E^\varphi \left(\frac{\bar{\lambda}^2}{\lambda^2 E^x} + \left\{ \frac{1}{E^x} - 2 \frac{\partial \ln \lambda^2}{\partial E^x} \right\} \frac{\sin^2 \bar{\lambda} K_\varphi}{\bar{\lambda}^2} \right) + 4K_x \frac{\sin 2\bar{\lambda} K_\varphi}{2\bar{\lambda}} \left(1 + \left(\frac{\bar{\lambda} (E^x)'}{2E^\varphi} \right)^2 \right) \right. \\ &\quad \left. + \left(\frac{(\{E^x\}')^2}{E^\varphi} \left(\frac{1}{2} \frac{\partial \ln \lambda^2}{\partial E^x} - \frac{1}{4E^x} \right) + \left(\frac{(E^x)'}{E^\varphi} \right)' \right) \cos^2 \bar{\lambda} K_\varphi \right] \end{aligned} \quad (\text{B1})$$

where the Hamiltonian is now periodic in K_φ and describes the same physical system because it is canonically related to the original constraint (A7). To preserve the structure of the hypersurface deformation algebra—albeit with modified structure functions—the constraint algebra requires that the Hamiltonian's dependence on K_x remains linear, which is related to the standard assumption that the diffeomorphism constraint maintains its classical expression. In LQG, one traditionally represents the Hamiltonian constraint

on the space of solutions that satisfy the diffeomorphism constraint. Similarly, the classical diffeomorphism constraint can be solved in order to eliminate K_x from the Hamiltonian constraint. The corresponding reduced phase space is obtained in combination with the partial gauge fixing choosing $q_{\theta\theta} = E^x$ in the areal gauge, $E^x = x^2$.

In this partially gauge-fixed system, the Hamiltonian constraint can be written as

$$\begin{aligned} \tilde{H}_{\text{grav}} &= -\frac{\bar{\lambda} E^\varphi}{\lambda 2x} \left(\frac{\bar{\lambda}^2}{\lambda^2} + (1 - x(\ln \lambda^2)') \frac{\sin^2 \bar{\lambda} K_\varphi}{\bar{\lambda}^2} \right) + \frac{\chi}{2\lambda \bar{\lambda}} E^\varphi \left(1 + \left(\frac{\bar{\lambda} x}{E^\varphi} \right)^2 \right) (\cos^2 \bar{\lambda} K_\varphi)' \\ &+ \chi \frac{\bar{\lambda}}{\lambda} \frac{x}{2E^\varphi} \left(\left\{ 1 - x(\ln \lambda^2)' + 2E^\varphi \left(\frac{x}{E^\varphi} \right)' \right\} \right) \cos^2 \bar{\lambda} K_\varphi. \end{aligned} \quad (\text{B2})$$

Given the fact that in LQG one has well-defined operators for $\hat{E}^\varphi$ and $\exp(i\bar{\lambda} K_\varphi)$, the Hamiltonian constraint (with a suitable discretization of spatial derivatives) can be represented as a well-defined operator on a Hilbert space in which each $E^\varphi(x)$ is Bohr compactified as in loop quantizations. The partial gauge fixing used here is acceptable

because it is imposed after demonstrating covariance of the full system of constraints. We also note the importance of systematic considerations of canonical transformations, which provide the equivalence of the Hamiltonian constraint (A7) with a decreasing holonomy function, and the Hamiltonian constraint (B1) strictly periodic in K_φ [3].