

Quantum geometry and the Schwarzschild singularity

Abhay Ashtekar^{1,2} and Martin Bojowald^{2,1}

1. Institute for Gravitational Physics and Geometry,

Physics Department, Penn State, University Park, PA 16802, USA

2. Max-Planck-Institut für Gravitationsphysik, Albert-Einstein-Institut,

Am Mühlenberg 1, D-14476 Potsdam, Germany

Abstract

In homogeneous cosmologies, quantum geometry effects lead to a resolution of the classical singularity without having to invoke special boundary conditions at the singularity or introduce ad-hoc elements such as unphysical matter. The same effects are shown to lead to a resolution of the Schwarzschild singularity. The resulting quantum extension of space-time is likely to have significant implications to the black hole evaporation process. Similarities and differences with the situation in quantum geometrodynamics are pointed out.

PACS numbers: 04.60.Pp, 04.70.Dy, 04.60.Nc, 03.65.Sq

I. INTRODUCTION

General relativity provides a subtle and powerful interplay between gravity and geometry, thereby opening numerous possibilities for novel phenomena. Among the most spectacular of the resulting conceptual advances are the predictions that the universe began with a big bang and massive stars can end their lives as black holes. In both cases, one encounters singularities. Space-time of general relativity literally ends and classical physics comes to a halt.

However, general relativity is incomplete because it ignores quantum effects. It is widely believed that quantum gravity effects become significant in the high curvature regions that develop before singularities are formed. These are likely to significantly change the space-time structure, making the predictions of general relativity unreliable. Hence, real physics need not stop at the big bang and black hole singularities. While the classical space-time does end there, quantum space-time may well extend beyond.

Attempts at extending physics beyond the big-bang singularity date back at least to the seventies when mini-superspaces were introduced and quantum cosmology was born (see, e.g., [1]). These investigations sparked off new developments in different directions ranging from foundations of quantum mechanics to the development of WKB methods to test semi-classicality in quantum cosmology, to the introduction of novel Euclidean methods to calculate the appropriate wave function of the universe. This deepened our understanding of issues related to the quantum physics of the universe as a whole [2]. However, in these approaches the singularity was not generically resolved. A quantum extension of space-time either required the introduction of new principles [3], or ad-hoc assumptions, such as existence of matter violating energy conditions already at the classical level, or of external clocks that remain insensitive to the infinite curvature encountered at singularities.

Emergence of quantum Riemannian geometry and the associated mathematical techniques [4, 5, 6] have provided a new approach to revisit the issue of quantum extensions. Within the mini-superspaces used in quantum cosmology, the big-bang singularity could be resolved *without* having to introduce external boundary conditions at the big-bang, or matter/clocks with unphysical properties [7, 8, 9]. In the homogeneous, isotropic mini-superspace coupled to a massless scalar field, in particular, an exhaustive analysis can be carried out [10]. At the analytic level, one can introduce an appropriate inner product on

physical states, define Dirac observables and, using the two, construct semi-classical states. One can then use numerical methods to examine the nature of the quantum space-time. If one begins with a semi-classical state at large late times (‘now’) and evolves it back in time, it remains semi-classical till one encounters the deep Planck regime near the classical singularity. In this regime the quantum geometry effects dominate. However, the state becomes semi-classical again on the other side; the deep Planck region serves as a quantum bridge between two large, classical space-times. The model is too simple to be applied reliably to the actual universe we live in. However at a conceptual and mathematical level it demonstrates the new possibilities, systematically made available by quantum geometry.

It is then natural to ask if a similar resolution of singularities occurs also in the context of black holes. Now, there is an extensive literature on the nature of black hole singularities in the classical theory. In particular, detailed mathematical analysis of spherical collapse of uncharged matter was carried out by Christodoulou, Dafermos and others. It has led to the expectation that, for black holes formed through gravitational collapse, the singularity would be generically space-like (see, e.g., [11] and references therein). Therefore, it is natural to focus on this case first and ask for the nature of quantum space-time that would result via their resolution.¹ If the singularity is resolved and if the quantum geometry in the deep Planck regime again serves as a bridge to a large classical region beyond, there would be no information loss in the black hole formation and evaporation process [14] (see also [15, 16]). Pure states in the distant past would evolve to pure states in the distant future and one would have a space-time description of the entire process in a quantum mechanical setting. Thus, a detailed analysis of the fate of black hole singularities is of considerable importance also for the fundamental issue of whether the standard unitary evolution of quantum physics has to be modified in the setting of black holes.

The purpose of this note is to initiate this investigation using quantum geometry, in the setting of connection dynamics. We will focus just on the mini-superspace that is appropriate for describing the geometry interior to the horizon of a Schwarzschild black

¹ Indeed, one does not expect quantum gravity effects to resolve all singularities. If, for example, they led to a resolution of the time-like singularity of the negative energy Schwarzschild space-time, energy would be unbounded below in quantum gravity and the theory would have unphysical features [12]. Rather, one expects that there would be *no* physical states in quantum gravity that would resemble negative energy Schwarzschild geometry at large distances, whence the issue of ‘resolution’ would simply not arise.

hole. Since this mini-superspace is spatially homogeneous —of Kantowski-Sachs type— one can take over the techniques that have been developed in the setting of homogeneous cosmologies [8]. Our main result is that the quantum scalar constraint is indeed such that the singularity is resolved. The salient features of our analysis are as follows. First, we use a *self-adjoint* Hamiltonian constraint. Therefore, our analysis can serve as the point of departure for the construction of the physical Hilbert space either via the group averaging procedure [13] *or* via deparametrization of the theory as in [10]. Second, we spell out the symmetry reduction procedure that leads to the Kantowski-Sachs mini-superspace *within connection dynamics*. We will see in sections III and IV that the structure of the resulting phase space is important to the issue of singularity resolution.² In particular, this space is an extension of the phase space used in geometrodynamics in that one allows for degenerate triads (and hence 3-metrics). Thanks to this enlargement, points of the reduced phase space corresponding to the singularity do not constitute a boundary, whence the support of the wave function ‘beyond singularity’ can be interpreted geometrically. However, we emphasize that the results of this note constitute only initial steps for a more complete theory, e.g., along the lines of [10].

Our discussion is organized as follows. In section II we discuss the classical theory of the Kantowski-Sachs mini-superspace using connection-dynamics. In section III we analyze the kinematics of its quantum theory and in section IV we present quantum dynamics. Section V provides a brief summary and directions for future work.

II. CLASSICAL THEORY

This section is divided into two parts. In the first, we carry out the Kantowski-Sachs symmetry reduction of vacuum general relativity and in the second we discuss the structure of the resulting phase space.

² A quantization of the Kantowski-Sachs minisuperspace is available in the geometrodynamical framework, [17]. Unfortunately, it does not shed light on singularity resolution because the location of the singularity in the minisuperspace was misidentified. In the notation used in [17], the singularity lies at $(a = \infty, b = 0)$ —not at $(a = 0, b = 0)$ as assumed there. Therefore, neither is the inverse volume operator (17) of [17] bounded or well-defined at the singularity nor does the discrete equation (23) of [17] enable one to evolve across the singularity.

A. Symmetry reduction

The portion of the space-time interior to the horizon of a Schwarzschild black hole can be naturally foliated by 3-manifolds M with topology $\mathbb{R} \times \mathbb{S}^2$ such that the geometry on each slice is invariant under the Kantowski–Sachs symmetry group $G = \mathbb{R} \times \text{SO}(3)$. In the discussion of this model we closely follow [9] where the mathematical structure of the simpler, isotropic model is studied in detail.

As in any homogeneous situation, M acquires an equivalence class of positive definite metrics related by an overall constant. We fix one such metric ${}^oq_{ab}$ as well as an orthonormal triad ${}^oe_i^a$ and co-triad ${}^o\omega_a^i$, compatible with it. With the given symmetry, we choose the metric on the orbits of the $\text{SO}(3)$ action to be the unit 2-sphere metric, $d\vartheta^2 + \sin^2\vartheta d\varphi^2$ in polar coordinates ϑ, φ . Locally, the corresponding co-triad elements $d\vartheta$, and $\sin\vartheta d\varphi$ can be completed to an orthonormal co-triad by a third form α . Thanks to the Maurer–Cartan relations for the symmetry group, α must be closed and therefore exact on the given topology, whence it defines a third adapted coordinate x via $\alpha = dx$. In these coordinates the background metric has line element

$$ds_o^2 = dx^2 + d\vartheta^2 + \sin^2\vartheta d\varphi^2$$

and determinant ${}^oq = \sin^2\vartheta$.

In connection dynamics, the canonically conjugate pair consists of fields (A_a^i, E_i^a) of an $\text{SU}(2)$ connection 1-form A_a^i and a (possibly degenerate) triad E_i^a of density weight 1 on M . This pair is invariant under the symmetry group if it satisfies:

$$\mathcal{L}_\xi A = D\Lambda \quad \text{and} \quad \mathcal{L}_\xi E = [E, \Lambda] \tag{1}$$

for any symmetry vector field ξ^a and some generator Λ^i of local, $\text{SU}(2)$ gauge transformations (which may depend on ξ^a). One can verify that any such symmetric pair is gauge equivalent

to fields of the form:³

$$A = \tilde{c}\tau_3 dx + (\tilde{a}\tau_1 + \tilde{b}\tau_2)d\vartheta + (-\tilde{b}\tau_1 + \tilde{a}\tau_2) \sin \vartheta d\varphi + \tau_3 \cos \vartheta d\varphi \quad (2)$$

$$E = \tilde{p}_c \tau_3 \sin \vartheta \frac{\partial}{\partial x} + (\tilde{p}_a \tau_1 + \tilde{p}_b \tau_2) \sin \vartheta \frac{\partial}{\partial \vartheta} + (-\tilde{p}_b \tau_1 + \tilde{p}_a \tau_2) \frac{\partial}{\partial \varphi}. \quad (3)$$

where $\tilde{a}, \tilde{b}, \tilde{c}$ are real constants and τ^j is the standard basis in $\mathfrak{su}(2)$, satisfying $[\tau_i, \tau_j] = \epsilon_{ij}^k \tau_k$. However, this does not exhaust the freedom to perform local $\text{SU}(2)$ gauge transformations entirely: There is still the freedom to perform a *global* $\text{SU}(2)$ transformation along τ_3 which rotates the ‘vectors’ (\tilde{a}, \tilde{b}) and $(\tilde{p}_a, \tilde{p}_b)$. We will return to this freedom in the next sub-section.

From the invariant triad density (3) we can derive the corresponding co-triad ω :

$$\omega = \omega_c \tau_3 dx + (\omega_a \tau_1 + \omega_b \tau_2) d\vartheta + (-\omega_b \tau_1 + \omega_a \tau_2) \sin \vartheta d\varphi \quad (4)$$

where

$$\omega_c = \frac{\text{sgn} \tilde{p}_c \sqrt{\tilde{p}_a^2 + \tilde{p}_b^2}}{\sqrt{|\tilde{p}_c|}}; \quad \omega_b = \frac{\sqrt{|\tilde{p}_c|} \tilde{p}_b}{\sqrt{\tilde{p}_a^2 + \tilde{p}_b^2}}; \quad \text{and } \omega_a = \frac{\sqrt{|\tilde{p}_c|} \tilde{p}_a}{\sqrt{\tilde{p}_a^2 + \tilde{p}_b^2}}. \quad (5)$$

The co-triad in turn determines its spin-connection, the $\mathfrak{su}(2)$ -valued 1-form Γ . As one might expect from the structure of the triad, Γ turns out to be the standard ‘magnetic monopole’ connection:

$$\Gamma = \cos \vartheta \tau_3 d\varphi. \quad (6)$$

Consequently, the $\mathfrak{su}(2)$ -valued extrinsic curvature 1-form K is given by:

$$\gamma K := A - \Gamma = \tilde{c}\tau_3 dx + (\tilde{a}\tau_1 + \tilde{b}\tau_2)d\vartheta + (-\tilde{b}\tau_1 + \tilde{a}\tau_2) \sin \vartheta d\varphi. \quad (7)$$

Finally the curvature F_{ab} of A_a can be easily computed. Its dual, the $\mathfrak{su}(2)$ -valued vector density $B^a := \frac{1}{2} \eta^{abc} F_{bc}$ has the same invariant form as E^a :

$$B = (\tilde{a}^2 + \tilde{b}^2 - 1) \tau_3 \sin \vartheta \frac{\partial}{\partial x} + (\tilde{a}\tau_1 + \tilde{b}\tau_2) \tilde{c} \sin \vartheta \frac{\partial}{\partial \vartheta} + (-\tilde{b}\tau_1 + \tilde{a}\tau_2) \tilde{c} \frac{\partial}{\partial \varphi} \quad (8)$$

These expressions will be used in sections III and IV.

³ More precisely, invariant connections [18, 19] carry a non-negative, integer-valued topological charge which we have set equal to one in the above expressions. (In general, it would multiply the last contribution to (2); see, e.g., [20].) Each value of this charge gives rise to an independent sector of invariant connections. However, only the sector used here allows non-degenerate triads invariant under (1) with *the same* local gauge transformation in A and E . That A in (2) satisfies (1) can be verified as follows. If we denote the three generators of the $\text{SO}(3)$ action on M by $X = \sin \varphi \partial_\vartheta + \cot \vartheta \cos \varphi \partial_\varphi$, $Y = -\cos \varphi \partial_\vartheta + \cot \vartheta \sin \varphi \partial_\varphi$, and $Z = \partial_\varphi$, and the generator of the \mathbb{R} -action by t^a , then $\mathcal{L}_X A = D(\cos \varphi / \sin \vartheta) \tau_3$, $\mathcal{L}_Y A = D(\sin \varphi / \sin \vartheta) \tau_3$, $\mathcal{L}_Z A = 0$ and $\mathcal{L}_t A = 0$. Invariance of E in (3) follows analogously, or by noting that $E_i^a \delta A_a^i$ defines a gauge invariant and (spatially) constant 1-form on phase space.

B. The reduced phase space

The symmetry reduction procedure led us to a 6-dimensional reduced phase space $\tilde{\Gamma}$ with coordinates $(\tilde{a}, \tilde{b}, \tilde{c}; \tilde{p}_a, \tilde{p}_b, \tilde{p}_c)$. Let us begin by computing the symplectic structure $\tilde{\Omega}$ on it. The basic idea, as in all symmetry reductions, is to pull-back the symplectic structure of the full theory to the symmetry reduced phase space. However, since the full symplectic structure involves an integral over M and since the fields of interest are homogeneous also in the non-compact \mathbb{R} -direction, as usual, we are led to consider only a finite interval \mathcal{I} in the \mathbb{R} direction. Let the length of this interval (w.r.t. the fiducial metric ${}^oq_{ab}$) be L_o . Then, the symplectic structure on the reduced phase space is given by:

$$\tilde{\Omega} = \frac{L_o}{2\gamma G} (2d\tilde{a} \wedge d\tilde{p}_a + 2d\tilde{b} \wedge d\tilde{p}_b + d\tilde{c} \wedge d\tilde{p}_c) \quad (9)$$

where γ is the Barbero–Immirzi parameter and G the gravitational constant. (The volume of the \mathbb{S}^2 orbits of the $\text{SO}(3)$ action, defined by the fiducial metric, is the standard one, 4π .)

Because of the form (2) and (3) of invariant connections and triads, the vector constraint is automatically satisfied. However, because of the residual, global $\text{SU}(2)$ gauge freedom mentioned in section II A, the Gauss constraint is not. Inserting the invariant connection and triad into the Gauss constraint we obtain

$$\mathcal{C}_{\text{Gauss}} = \tilde{a}\tilde{p}_b - \tilde{b}\tilde{p}_a = 0 \quad (10)$$

which generates simultaneous rotations of the pairs (\tilde{a}, \tilde{b}) and $(\tilde{p}_a, \tilde{p}_b)$. Thus, only the ‘norms’ $\sqrt{\tilde{a}^2 + \tilde{b}^2}$, $\sqrt{\tilde{p}_a^2 + \tilde{p}_b^2}$ and the ‘scalar product’ $(\tilde{a}\tilde{p}_a + \tilde{b}\tilde{p}_b)$ are gauge invariant. We will fix this gauge freedom as follows. If $(\tilde{a}, \tilde{b}) = 0$, we rotate the triad components such that $\tilde{p}_a = 0$. Otherwise we rotate (\tilde{a}, \tilde{b}) such that $\tilde{a} = 0$ which implies $\tilde{b} \neq 0$. Then, $\mathcal{C}_{\text{Gauss}} = 0$ implies $\tilde{p}_a = 0$. There is still a residual gauge freedom $\Pi_b: (b, p_b) \rightarrow (-b, -p_b)$ which changes the signs of b and p_b simultaneously. This is just the parity transformation in the p_b variable. One can either retain this freedom and ensure that all the final constructions are invariant with respect to this ‘ b -reflection’ or eliminate it by a gauge choice such as $\tilde{p}_b \geq 0$. This gauge choice turns out to be particularly convenient for classical dynamics. In the quantum theory, on the other hand, it is more natural to retain the freedom at first and then ask that physical states be invariant under the parity operator $\hat{\Pi}_b$ implementing this transformation. Therefore, we will allow both possibilities. In either case the 4-dimensional phase space

carries a global chart $(\tilde{b}, \tilde{c}; \tilde{p}_b, \tilde{p}_c)$ (where $\tilde{p}_b \geq 0$ if the ‘ b -parity gauge’ is fixed.) These coordinates are subject to the scalar or the Hamiltonian constraint, discussed below.

We first note that these variables and, because of its explicit dependence on L_o , the symplectic structure $\tilde{\Omega}$ depend on the fiducial metric. It is convenient to remove this dependence by rescaling the variables in a manner that is motivated by their scaling properties,

$$(b, c) := (\tilde{b}, L_o \tilde{c}) \quad , \quad (p_b, p_c) := (L_o \tilde{p}_b, \tilde{p}_c) . \quad (11)$$

We now have a one-to-one parametrization of the gauge fixed phase space Γ by (b, p_b, c, p_c) with symplectic structure

$$\Omega = \frac{1}{2\gamma G} (2db \wedge dp_b + dc \wedge dp_c) . \quad (12)$$

For later purposes let us express the volume of our elementary cell $\mathcal{I} \times \mathbb{S}^2$ and areas bounded by our preferred family of curves as functions on the reduced phase space. The volume is given by:

$$V = \int d^3x \sqrt{|\det E|} = 4\pi L_o \sqrt{|\tilde{p}_c|} |\tilde{p}_b| = 4\pi \sqrt{|p_c|} |p_b| . \quad (13)$$

The three surfaces $S_{x,\varphi}$, $S_{x,\vartheta}$ and $S_{\vartheta,\varphi}$ of interest are respectively bounded by the interval \mathcal{I} and the equator, \mathcal{I} and a great circle along a longitude, and the equator and a longitude (so that $S_{\vartheta,\varphi}$ forms a quarter of the sphere \mathbb{S}^2). Their areas are given by:

$$A_{x,\vartheta} = 2\pi |p_b|, \quad A_{x,\varphi} = 2\pi |p_b|, \quad \text{and} \quad A_{\vartheta,\varphi} = \pi |p_c| . \quad (14)$$

Finally, let us consider the Hamiltonian constraint and classical dynamics. On the reduced phase space Γ , the constraint functional of the full theory

$$\mathcal{C}_{\text{Ham}} = \int d^3x N e^{-1} [\epsilon_{ijk} F_{ab}^i E_j^a E_k^b - 2(1 + \gamma^2) K_{[a}^i K_{b]}^j E_i^a E_j^b] , \quad (15)$$

where $e := \sqrt{|\det E|} \text{sgn}(\det E)$, reduces to

$$\mathcal{C}_{\text{Ham}} = -\frac{8\pi N}{\gamma^2} \frac{\text{sgn} p_c}{\sqrt{|p_c|} p_b} [(b^2 + \gamma^2) p_b^2 + 2c p_c b p_b] , \quad (16)$$

where, as is usual in the homogeneous models, we have chosen a constant lapse function N . To simplify the equations of motion, it is convenient⁴ to choose $N = \frac{\gamma \text{sgn} p_c \sqrt{|p_c|}}{16\pi G b}$. Then

⁴ With this choice, the points $b = 0$ or $p_c = 0$ have to be excised in the discussion of dynamics. The explicit form of the solutions below shows that this is not overly restrictive since on classical solutions these points correspond to the singularity or the horizon.

the Hamiltonian constraint becomes $\mathcal{C}_{\text{Ham}} := -(2\gamma G)^{-1}[(b^2 + \gamma^2)p_b/b + 2cp_c]$, yielding the following equations of motion:

$$\begin{aligned}\dot{b} &= -\frac{1}{2}(b + \gamma^2 b^{-1}), & \dot{p}_b &= \frac{1}{2}(p_b - \gamma^2 b^{-2} p_b) \\ \dot{c} &= -2c, & \dot{p}_c &= 2p_c,\end{aligned}$$

where the ‘dot’ denotes derivative with respect to the affine parameter of the Hamiltonian vector field. The equations for (c, p_c) can easily be solved by $c(T) = c_0 e^{-2T}$ and $p_c(T) = p_c^{(0)} e^{2T}$. For the other components we obtain $b(T) = \pm \gamma \sqrt{e^{-(T-T_0)} - 1}$ for b and, using this result, $p_b(T) = p_b^{(0)} \sqrt{e^{T+T_0} - e^{2T}}$. If we introduce the new time parameter $t := e^T$ and the constant $m := \frac{1}{2} e^{T_0}$, this gives

$$b(t) = \pm \gamma \sqrt{(2m - t)/t}, \quad p_b(t) = p_b^{(0)} \sqrt{t(2m - t)} \quad (17)$$

$$c(t) = \mp \gamma m p_b^{(0)} t^{-2}, \quad p_c(t) = \pm t^2 \quad (18)$$

where we have fixed the multiplicative constant in the expressions of $p_c(t)$ and $c(t)$, respectively, by requiring that $|p_c|$ be the geometric radius of the 2-sphere orbits of the $\text{SO}(3)$ symmetry and by using the Hamiltonian constraint. Projections to the $p_b - p_c$ plane of typical trajectories are shown in figure 1. It is clear from the figure that p_c can be taken to be the ‘internal time parameter’. In section IV A this interpretation will let us regard the quantum Hamiltonian constraint as an ‘evolution equation’ with respect to an internal time parameter defined by the eigenvalues of the operator \hat{p}_c .

Note that we have a 2-parameter family of solutions, labeled by m and $p_b^{(0)}$, as one would expect from the fact that the reduced phase space is 2-dimensional. However, in a space-time description we expect only a 1-parameter family, labeled by m . This discrepancy can be traced back to the standard tension between the Hamiltonian and space-time notions of gauge (for a detailed discussion in the context of Bianchi models, see [21], and in spherically symmetric models, [22, 23]). In the Hamiltonian description, rescalings of $p_b^{(0)}$ do not correspond to gauge because they are not generated by any constraint. In the space-time description, on the other hand, they can be absorbed by a rescaling of the coordinate x . Therefore, to make contact with the space-time description, let us fix $p_b^{(0)} = 1$. Then, the above solution to the evolution equations defines a 1-parameter family of 3-metrics:

$$q_{ab}(t) = \left(\frac{2m}{t} - 1\right) \nabla_a x \nabla_b x + t^2 (\nabla_a \vartheta \nabla_b \vartheta + \sin^2 \vartheta \nabla_a \varphi \nabla_b \varphi) \quad (19)$$

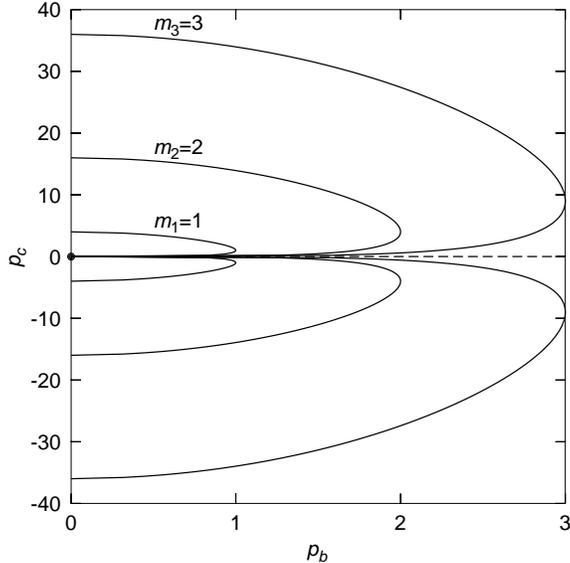


FIG. 1: Dynamical trajectories in the $p_b - p_c$ plane. Each trajectory reaches the maximum value m of p_b and meets the $p_c = 0$ axis only at the point where p_b also vanishes (i.e., at the ‘origin’). Solutions related by the reflection-symmetry $p_c \rightarrow -p_c$ define the same metric but carry triads of opposite orientation. For simplicity the ‘ b -parity gauge’ is fixed by requiring $p_b \geq 0$.

defining precisely the Schwarzschild solution of mass m , interior to the horizon. It will be useful to note that, in this solution $p_c > 0$, $p_b = 0$ at the horizon ($t = 2m$) and $p_c = 0$ and b, c diverge at the singularity.

We will see in the next section that it is convenient to regard the space \mathcal{M} spanned by p_b and p_c as the mini-superspace on which quantum wave functions are defined. (In that discussion, we will not fix the ‘ b -parity gauge’; so p_c and p_b take values on the entire real line.) Let us therefore interpret various regions in \mathcal{M} . The triad is non-degenerate everywhere except when $p_b = 0$ or $p_c = 0$. However, the two degeneracies are of very different geometric and physical origin. Geometrically (i.e., independent of classical dynamics) $p_c = 0$ separates two regions with triads of opposite orientations, given by $\text{sgn det } E = \text{sgn}(p_c p_b^2)$. On the other hand, as (4) (together with our gauge choice which ensures $a = p_a = 0$) shows, the orientation would not have changed even if we had extended \mathcal{M} across $p_b = 0$, allowing for negative values for p_b . More importantly, while the co-triad (4) remains smooth and only becomes degenerate along $p_b = 0$, it diverges along $p_c = 0$, signaling the classical singularity. This suggests that the line $p_b = 0$ corresponds to the horizon and $p_c = 0$ to

the singularity. This is confirmed from physical considerations involving classical dynamics. Classical solutions can be smoothly extended to the line $p_b = 0$ which is reached for $t = 2m$. The curvature B remains well-defined there (see (8)). The other line, $p_c = 0$, is very different. Along classical solutions, the line $p_c = 0$ is approached as $t \rightarrow 0$, whence the components (and invariants) of curvature B also diverge. This line corresponds to the singularity and cannot be crossed by any classical dynamical trajectory. Finally, note that since the line $p_c = 0$ separates the two regions of \mathcal{M} , where only the triad orientation is reversed, the two get identified in geometrodynamics. Therefore the singularity would lie at a boundary of the mini-superspace in geometrodynamics, making the geometrical meaning of ‘evolution across the singularity’ obscure. Furthermore, since the Arnowitt–Deser–Misner (ADM) variables are based on covariant metrics (rather than contravariant triad densities), as in other homogeneous models [8], one of the ADM variables becomes *infinite* at the classical singularity. This further complicates a discussion of the singularity structure and its resolution.

III. QUANTUM KINEMATICS

This section is divided into two parts. In the first we introduce the basic quantization procedure and in the second we discuss quantum geometry with emphasis on the classical singularity.

A. Basics

As in the isotropic model [9], we will follow the general procedure used in the full theory [4, 5, 6]. Thus the elementary configuration variables will be given by holonomies along curves in M and the momenta by fluxes of triads along 2-surfaces in M . However, because of symmetry reduction, one need not consider all (piecewise analytic) curves or surfaces; a judiciously chosen, smaller subset suffices to obtain a set of functions that is sufficiently large to separate points of the reduced phase space.

Let us begin with the holonomies. We will restrict ourselves to three sets of curves: those along the \mathbb{R} direction of M with oriented length τL_o ; those along the equator of \mathbb{S}^2 with oriented length μ ; and those along the longitudes of \mathbb{S}^2 also with oriented length μ , where

all lengths and orientations are defined using the fiducial triad. (Thus τ, μ are positive if the tangent to the respective curves are parallel to the triad vectors and negative if they are anti-parallel.) Holonomies along these curves suffice to completely determine any invariant connection (2).⁵

$$h_x^{(\tau)}(A) = \exp \int_0^{\tau L_o} dx \tilde{c} \tau_3 = \cos \frac{\tau c}{2} + 2\tau_3 \sin \frac{\tau c}{2} \quad (20)$$

$$h_\varphi^{(\mu)}(A) = \exp - \int_0^\mu d\varphi \tilde{b} \tau_1 = \cos \frac{\mu b}{2} - 2\tau_1 \sin \frac{\mu b}{2} \quad (21)$$

$$h_\vartheta^{(\mu)}(A) = \exp \int_0^\mu d\vartheta \tilde{b} \tau_2 = \cos \frac{\mu b}{2} + 2\tau_2 \sin \frac{\mu b}{2}. \quad (22)$$

Matrix elements of these holonomies are functions of the reduced connection and constitute our configuration variables. Elements of the algebra they generate are almost periodic functions of b and c of the form $f(b, c) = \sum_{\mu, \tau} f_{\mu, \tau} \exp \frac{i}{2} (\mu b + \tau c)$, where $f_{\mu, \tau} \in \mathbb{C}$, $\mu, \tau \in \mathbb{R}$ and the sum extends over a finite set. This algebra is the Kantowski-Sachs analog of the algebra of cylindrical functions in the full theory [4, 5, 6]. We will therefore denote it by Cyl_{KS} . Consider the C^* algebra obtained by completing Cyl_{KS} using the sup norm. The quantum configuration space $\bar{\mathcal{A}}_{\text{KS}}$ is the Gel'fand spectrum of this algebra. From the structure of the algebra, it follows that the spectrum is naturally isomorphic to the Bohr compactification $\bar{\mathbb{R}}_{\text{Bohr}}^2$ of the Abelian group \mathbb{R}^2 [25, 26]. (Recall that in the isotropic case [9], the quantum configuration space is isomorphic to $\bar{\mathbb{R}}_{\text{Bohr}}$.) The Hilbert space $\tilde{\mathcal{H}}$ is obtained by the Cauchy completion of Cyl_{KS} with respect to the natural Haar measure μ_o on the Abelian group $\bar{\mathbb{R}}_{\text{Bohr}}^2$; $\tilde{\mathcal{H}} = L^2(\bar{\mathbb{R}}_{\text{Bohr}}^2, d\mu_o)$. Employing the standard bra-ket notation, we can define a basis $|\mu, \tau\rangle$ in $\tilde{\mathcal{H}}$ via:

$$\langle b, c | \mu, \tau \rangle = e^{\frac{i}{2} (\mu b + \tau c)} \quad \mu, \tau \in \mathbb{R}. \quad (23)$$

This is an orthonormal basis:

$$\langle \mu', \tau' | \mu, \tau \rangle = \delta_{\mu' \mu} \delta_{\tau' \tau} \quad (24)$$

where, on the right side, we have the Kronecker symbol, rather than the Dirac delta distribution. Thus, each basis vector is normalizable and has unit norm.

⁵ In the fully gauge fixed setting, we could omit h_ϑ since it is related to h_φ by conjugation with $\exp \tau_3$. However, we will use a more ‘democratic’ approach which is applicable also when a is not made to vanish by gauge fixing.

As one would expect, the configuration —i.e., holonomy— operators $\hat{h}_x^{(\tau)}$, $\hat{h}_\varphi^{(\mu)}$, $\hat{h}_\theta^{(\mu)}$ operate on $\tilde{\mathcal{H}}$ by multiplication. However, as in the full theory, these operators are not required to be weakly continuous in parameters μ and τ , whence there are no operators \hat{b}, \hat{c} corresponding to the connection itself. consequently, although we have only a finite number of degrees of freedom, the von-Neumann uniqueness theorem is inapplicable and this quantum theory is inequivalent to a standard ‘Schrödinger quantization’ (for a further discussion, see [25, 26]). We will see that, as in the isotropic model [9], this inequivalence has important consequences.

For the momentum operators we consider fluxes of triads along preferred 2-surfaces. Apart from fixed kinematical factors, they are given by components p_b, p_c of triads which, in view of the symplectic structure (12), are represented by operators⁶

$$\hat{p}_b = -i\gamma\ell_{\text{Pl}}^2 \frac{\partial}{\partial b}, \quad \hat{p}_c = -2i\gamma\ell_{\text{Pl}}^2 \frac{\partial}{\partial c}. \quad (25)$$

Their eigenstates are the basis states (23),

$$\hat{p}_b|\mu, \tau\rangle = \frac{1}{2}\gamma\ell_{\text{Pl}}^2\mu|\mu, \tau\rangle, \quad \hat{p}_c|\mu, \tau\rangle = \gamma\ell_{\text{Pl}}^2\tau|\mu, \tau\rangle. \quad (26)$$

However, we still have to incorporate the residual gauge freedom which corresponds to a parity reflection in the b degree of freedom. Therefore, only those states in $\tilde{\mathcal{H}}$ which are invariant under the parity operator $\hat{\Pi}_b: |\mu, \tau\rangle \rightarrow |-\mu, \tau\rangle$ can belong to the kinematical Hilbert space \mathcal{H} . A basis in \mathcal{H} is thus given by:

$$\frac{1}{\sqrt{2}} [|\mu, \tau\rangle + |-\mu, \tau\rangle] \quad (27)$$

Finally, we express the volume operator in terms of triad operators \hat{p}_b, \hat{p}_c . Recall that the region of M under consideration is of the type $\mathcal{I} \times \mathbb{S}^2$, where \mathcal{I} has length L_o with respect to the fiducial metric. From the classical expression (13) of the volume V of this region, it follows that the operator \hat{V} is given by $\hat{V} = 4\pi|\hat{p}_b|\sqrt{|\hat{p}_c|}$. It is diagonal in our $|\mu, \tau\rangle$ basis and the eigenvalues are:

$$V_{\mu\tau} = 2\pi\gamma^{3/2}|\mu|\sqrt{|\tau|}\ell_{\text{Pl}}^3. \quad (28)$$

As in the isotropic case, the volume of our cell $\mathcal{I} \times \mathbb{S}^2$ with respect to the fiducial metric can be large, but its volume in the ‘elementary’ state $\exp \frac{i}{2}(b+c)$, is of Planck size.

⁶ In this paper, we use the standard quantum gravity convention, $\ell_{\text{Pl}}^2 = G\hbar$. Unfortunately, this is different from the convention $\ell_{\text{Pl}}^2 = 8\pi G\hbar$ used in most of the loop quantum cosmology literature. Hence care should be exercised while comparing detailed numerical factors.

B. Quantum geometry

As in the isotropic case the triad operators \hat{p}_b, \hat{p}_c commute whence, in contrast to the full theory, one can construct the triad representation. In this sub-section we will study the quantum Riemannian geometry using this representation. In the next sub-section we will see that the representation is also convenient for dynamics, for it suggests an intuitive interpretation for the action of the Hamiltonian constraint.

Let us then expand a general state in terms of eigenstates of the triad operators, $|\psi\rangle = \sum_{\mu\tau} \psi_{\mu\tau} |\mu, \tau\rangle$, and use the coefficients $\psi_{\mu\tau}$ to represent the state. This wave function $\psi_{\mu\tau}$ is supported on the mini-superspace \mathcal{M} coordinatized by p_b and p_c . In the classical theory, already at the kinematical level, the lines $p_b = 0$ and $p_c = 0$ are special: the co-triad (4) becomes degenerate along $p_b = 0$ (and the line represents the horizon) while it diverges on the line $p_c = 0$ (and this line represents the singularity). It is then natural to examine wave functions $\psi_{\mu\tau}$ which are supported on lines $\mu = 0$ and $\tau = 0$ and ask, already at the kinematical level, for the nature of quantum geometry they represent.

Let us consider the co-triad component $\omega_c \equiv \text{sgn}(p_c) |p_b|/\sqrt{|p_c|}$ along dx (see (4)). Since it involves an inverse power of p_c and since there exist normalizable kets $|\mu, \tau = 0\rangle$ of the operator \hat{p}_c with zero eigenvalue, the naive operator obtained by replacing p_b and p_c by corresponding operators is not even densely defined, let alone self-adjoint. To define $\hat{\omega}_c$, therefore, a new strategy is needed. We will follow what is by now a standard procedure, adapted from Thiemann's analysis in the full theory [24]. The first step is to express ω_c in terms of the elementary variables which do have unambiguous quantum analogs —holonomies and positive powers of p_b, p_c — and Poisson brackets between them. On the classical phase space, we have the exact equality:

$$\omega_c = \frac{1}{2\pi\gamma G} \text{Tr} \left(\tau_3 h_x \{h_x^{-1}, V\} \right) \quad (29)$$

where h_x is the holonomy along the interval \mathcal{I} , i.e., along the edge of length L_o (with respect to the fiducial metric) we fixed in our construction of the phase space. Then, replacing h_x and V by their unambiguous quantum analogs and the Poisson bracket by $1/i\hbar$ times the commutator, we obtain

$$\begin{aligned} \hat{\omega}_c &= -\frac{i}{2\pi\gamma\ell_{\text{Pl}}^2} \text{Tr} \left(\tau_3 \hat{h}_x [\hat{h}_x^{-1}, \hat{V}] \right) \\ &= -\frac{i}{2\pi\gamma\ell_{\text{Pl}}^2} \left(\sin \frac{c}{2} \hat{V} \cos \frac{c}{2} - \cos \frac{c}{2} \hat{V} \sin \frac{c}{2} \right). \end{aligned} \quad (30)$$

It turns out that this operator is diagonalized by our basis $|\mu, \tau\rangle$ of (23),

$$\hat{\omega}_c|\mu, \tau\rangle = \frac{1}{4\pi\gamma\ell_{\text{Pl}}^2}(V_{\mu, \tau+1} - V_{\mu, \tau-1})|\mu, \tau\rangle = \frac{\sqrt{\gamma}}{2}\ell_{\text{Pl}}|\mu|(\sqrt{|\tau+1|} - \sqrt{|\tau-1|})|\mu, \tau\rangle. \quad (31)$$

For large τ , far away from the classical singularity, these eigenvalues are very close to the classical expectation $|\omega_c| = |p_b|/\sqrt{|p_c|}$. This is even true for small μ , i.e. we do not need to be far from the horizon. Closer to the classical singularity $p_c = 0$, on the other hand, the behavior becomes significantly different from the classical one. At the classical singularity itself, where $|p_b|/\sqrt{|p_c|}$ diverges, the eigenvalue of $\hat{\omega}_c$ is zero.⁷

Remark: In the above definition of $\hat{\omega}_c$ we needed a holonomy in the x -direction. We chose to evaluate it using the interval \mathcal{I} of length L_o with respect to the fiducial metric. While this choice is natural because this interval appears already in the construction of the phase space, we could have used an interval of length $L_o\delta$ for some δ . Then the resulting operator would also have been bounded on the entire Hilbert space $\tilde{\mathcal{H}}$, and $|\mu, \tau = 0\rangle$ would again have been eigenvectors of the operator with zero eigenvalue. Thus, the qualitative properties of the operator are insensitive to δ . Furthermore, from general considerations one can argue [9] that δ should be of the order of 1. However, the value of the upper bound of the spectrum does depend on δ . Therefore, as in the isotropic model, the precise numerical coefficient in this bound should not be attributed physical significance.

Let us now turn to the second triad component $\omega_b = \text{sgn}p_b \sqrt{|p_c|}$. Since $\sqrt{|\tau|}$ is a well-defined function of τ on the entire spectrum of \hat{p}_c , we can quantize ω_b directly. This operator is again diagonal in our basis $|\mu, \tau\rangle$ and its eigenvalues are given by

$$\hat{\omega}_b|\mu, \tau\rangle = \sqrt{\gamma}\ell_{\text{Pl}}\text{sgn}(\mu)\sqrt{|\tau|}|\mu, \tau\rangle. \quad (32)$$

Since the operator is quantized directly, the eigenvalues are just the ones one would expect from the classical expression (4) of ω_b .

The properties of co-triad operators suggest that the lines $\mu = 0$ and $\tau = 0$ have very different features also in quantum geometry. No quantum effects are manifest at $\mu = 0$ which corresponds to the classical horizon. Near and on the line $\tau = 0$ which represents the

⁷ Exact classical identities such as (29) for objects containing inverse powers are available only in homogeneous models. More generally, the bracket between holonomies and volume is more complicated [28]. These complications lead to additional correction terms and the resulting operator is no longer bounded on eigenspaces of \hat{p}_b .

singularity, on the other hand, quantum effects are large. In particular, they remove the classical singular behavior of the co-triad. We emphasize, however, that even in homogeneous models, the ultimate test as to whether or not a singularity persists upon quantization can only come from studying the dynamics. The key questions are: Is quantum dynamics well-defined and deterministic across the classical singularity? And, does this come about without significant quantum corrections to geometry near the horizon? Only quantum dynamics will tell us if the indication provided by the properties of the co-triad operators are borne out.

IV. QUANTUM DYNAMICS

This section is divided into three parts. In the first we present the strategy, in the second we obtain the quantum Hamiltonian constraint, and in the third we use it to discuss the consequent resolution of the classical singularity.

A. Strategy

To bring out the similarities and differences between the reduced model and the full theory, it is instructive to begin with the Hamiltonian constraint (15) of the full theory:

$$\mathcal{C}_{\text{Ham}} = \int d^3x N e^{-1} [\epsilon_{ijk} E^{ai} E^{bj} F_{ab}^k - 2(1 + \gamma^2) E_i^a E_j^b K_{[a}^i K_{b]}^j]. \quad (33)$$

We begin by noting some simplifications that arise because of spatial homogeneity. Recall that the connection A is related to the spin-connection Γ defined by the triad and the extrinsic curvature K via $A = \Gamma + \gamma K$. As remarked in section II A, symmetries of the model imply that Γ is the standard magnetic monopole connection; it does not depend on the phase space point under consideration. From its expression (6), it follows that its curvature Ω is given by:

$$\Omega = -\sin \vartheta \tau_3 d\vartheta \wedge d\varphi \quad (34)$$

Since it is a ‘c-number’ it can be trivially taken over to quantum theory. Now, the curvature F of the full connection A can be expanded out as:

$$F_{ab} = 2\partial_{[a} A_{b]} + [A_a, A_b] = \Omega_{ab} + 2\gamma\partial_{[a} K_{b]} + \gamma^2[K_a, K_b] + \gamma[\Gamma_a, K_b] - \gamma[\Gamma_b, K_a] \quad (35)$$

Using the expressions (6), (7) and (3) of Γ , K and E , it is straightforward to verify that:

$$\begin{aligned} [\Gamma, \gamma K] &= b\tau_2 \cos \vartheta d\vartheta \wedge d\varphi = -\gamma dK \\ \epsilon_{ijk}(\partial_{[a}K_{b]}^i) E_j^a E_k^b &= 0 \end{aligned} \quad (36)$$

These equations now imply that the integrand of the Hamiltonian constraint can be written simply as⁸

$$\mathcal{C}_{\text{Ham}}(x) = e^{-1} E^{ai} E^{bj} (\epsilon_{ijk} F_{ab}^k - 2(1 + \gamma^2) K_{[a}^i K_{b]}^j) = e^{-1} E^{ai} E^{bj} (\epsilon_{ijk} \Omega_{ab}^k - 2K_{[ai} K_{b]j}). \quad (37)$$

Since Ω is constant on the entire phase space, the non-trivial part of the Hamiltonian constraint is thus contained just in a term quadratic in extrinsic curvature.

To pass to the quantum theory, we have to express the right hand side in terms of holonomies and triads. There are two possible avenues.

i) Using (36), the extrinsic curvature terms in the constraint function $\mathcal{C}_{\text{Ham}}(x)$ can be written in terms of curvatures F and Ω so that we have:

$$\mathcal{C}_{\text{Ham}}(x) = \frac{1}{\gamma^2} \epsilon_{ijk} e^{-1} E^{ai} E^{bj} [(1 + \gamma^2) \Omega_{ab}^k - F_{ab}^k]. \quad (38)$$

The idea now is to use the discussion of co-triads in section IIIB to obtain the operator analog of $\epsilon_{ijk} e^{-1} E^{ai} E^{bj}$; carry Ω to quantum theory trivially and use holonomies to express the field strength F as in the full theory.

ii) Alternatively, since we have completely fixed the gauge freedom to perform internal $\text{SU}(2)$ rotations, we can regard K itself as a connection. It is a direct analog of the connection A in the spatially flat cosmologies where Γ vanishes. Therefore, we will denote its curvature by oF ; ${}^oF = dK + [K, K]$. Then, the constraint functional can also be expressed as:

$$\mathcal{C}_{\text{Ham}}(x) = \frac{1}{\gamma^2} \epsilon_{ijk} e^{-1} E^{ai} E^{bj} [\gamma^2 \Omega_{ab}^k - {}^oF_{ab}^k] \quad (39)$$

Now, one can again use the discussion of co-triads in section IIIB to obtain the operator analog of $\epsilon_{ijk} e^{-1} E^{ai} E^{bj}$; carry Ω to quantum theory trivially and use holonomies to express the field strength oF . However, now holonomies have to be constructed using the connection K rather than A . Since our Hilbert space \mathcal{H} is built from holonomies of A , at first it

⁸ This simplification has a natural origin. Because of (36), the right hand side is just $\epsilon_{ijk} E^{ai} E^{bj} {}^+F_{ab}^k$, where ${}^+F$ is the curvature of the self-dual connection ${}^+A = \Gamma + iK$.

seems difficult to express holonomies of K as operators on \mathcal{H} . However, because of spatial homogeneity, we can calculate $\mathcal{C}_{\text{Ham}}(x)$ at any point. Let us choose a point on the equator and use holonomies along the three sets of curves introduced in (20)–(22). Using the expressions (2) and (7) of A and K , it follows that, along these curves the holonomies of A and K are equal! Therefore, the required holonomies of K can indeed be expressed as operators on \mathcal{H} .

Both these strategies are viable and, since they begin with *exact* expressions (38) and (39) of classical constraints, the difference in the resulting operators is just a quantization ambiguity. The first strategy is more natural in the sense that it does not require the introduction of a new connection. On the other hand, since the new connection K is the direct analog of the connection used in spatially flat cosmologies (in particular in [9]), to facilitate comparisons, the second strategy has been used in general, homogeneous models [8, 30]. For definiteness, we will use the second strategy in the main discussion and comment at the end on how the first strategy modifies the final constraint operator. The analysis of the resolution of the singularity can be carried out with either methods and the conclusion is the same.

B. Quantum Hamiltonian Constraint

Starting from the classical Hamiltonian constraint

$$\mathcal{C}_{\text{Ham}} = \frac{1}{\gamma^2} \int d^3x \epsilon_{ijk} e^{-1} E^{ai} E^{bj} [\gamma^2 \Omega_{ab}^k - {}^oF_{ab}^k], \quad (40)$$

where the integral is taken over the elementary cell, we wish to pass to the quantum operator. The overall strategy is the same as in [8, 9, 30] and the subtleties are discussed in detail in [9]. Therefore, here we will present only the main steps and comment on differences from the isotropic case treated in [9].

In the isotropic case, the elementary cell is a cube of length L_o with respect to the fiducial metric. To express the curvature of the connection and the co-triad operator in the Hamiltonian constraint, one uses holonomies along the curves of variable length $\mu_o L_o$ along the edges of this cell. For each value of μ_o one obtains a quantum constraint operator $\hat{\mathcal{C}}_{\text{Ham}}^{(\mu_o)}$. As discussed in [9], considerations from the full theory imply that μ_o should not be regarded as a regulator; rather, the μ_o -dependence of this operator should be regarded as a quantization ambiguity. This ambiguity could not be fixed within the reduced model itself.

But by making an appeal to results from the full theory, one can fix the value of μ_o using the minimum eigenvalue of the area operator in the full theory. (Recall that area enclosed by a loop enters while expressing curvature in terms of holonomies.)

In the present, Kantowski-Sachs model, the overall procedure is the same but minor adjustments are necessary because of lack of isotropy. Since \mathbb{S}^2 is compact, our elementary cell $\mathbb{S}^2 \times \mathcal{I}$ has a geometric edge only along the \mathcal{I} direction. As in section III A, we will supplement it with edges along a longitude and the equator of \mathbb{S}^2 . Holonomies along these three edges are then given by (20)–(22). Now consider curves of length $L_o\delta$ along \mathcal{I} and of length δ each along the equator and the longitude of \mathbb{S}^2 . Then, the co-triad function in the Hamiltonian (40) can be expressed as:

$$\epsilon_{ijk}\tau^i e^{-1}E^{aj}E^{bk} = -2(8\pi\gamma G\mathcal{L}_{(k)})^{-1}\epsilon^{abc}{}^o\omega_c^k h_k^{(\delta)}\{h_k^{(\delta)-1}, V\} \quad (41)$$

where on the right hand side the index k runs over x, ϑ, φ and is summed over; h_k is the holonomy along the edge k ; and $\mathcal{L}_x = L_o\delta$ and $\mathcal{L}_\vartheta = \mathcal{L}_\varphi = \delta$. Note that this is an exact equality for any choice of δ . Components of the curvature oF can also be expressed in terms of these holonomies. Set

$$h_{ij}^{(\delta)} = h_i^{(\delta)}h_j^{(\delta)}(h_i^{(\delta)})^{-1}(h_j^{(\delta)})^{-1} \quad (42)$$

where i, j, k run over x, ϑ, φ . Then, using (20)–(22) it is straightforward to verify that

$${}^oF_{ab}^i(x)\tau_i = \frac{{}^o\omega_a^i{}^o\omega_b^j}{\mathcal{A}_{(ij)}}(h_{ij}^{\mathcal{A}_{(ij)}} - 1) + O((b^2 + c^2)^{3/2}\sqrt{\mathcal{A}}) \quad (43)$$

where $\mathcal{A}_{x\vartheta} = \delta^2 L_o = \mathcal{A}_{x\varphi}$ and $\mathcal{A}_{\vartheta\varphi} = \delta^2$. For ${}^oF_{x,\vartheta}$ and ${}^oF_{x,\varphi}$, (43) is the standard geometric relation between holonomies around closed loops, the area they enclose and curvature. For ${}^oF_{\vartheta,\varphi}$, on the other hand, while (43) continues to hold, the standard geometrical interpretation is no longer available because the edges of length δ along the equator and a longitude fail to form a closed loop. Nonetheless, the standard relation is meaningfully extended because of spatial homogeneity, δ^2 playing the role of the ‘effective area’ $\mathcal{A}_{\vartheta,\varphi}$.

We now have all the ingredients, oF , Ω and the triad components, to rewrite the classical

constraint (40) in a way which is suitable for quantization. For the oF -term in (40) we obtain

$$\begin{aligned}
& -\gamma^{-2} \int d^3x \epsilon_{ijk} {}^oF_{ab}^i e^{-1} E^{aj} E^{bk} = 2\gamma^{-2} \int d^3x \text{Tr} ({}^oF_{ab}^i \tau_i \epsilon_{jkl} e^{-1} E^{aj} E^{bk} \tau^l) \\
& = -4(8\pi\gamma^3 G)^{-1} \int d^3x \det {}^o\omega \text{Tr} \left(\epsilon^{ijk} \mathcal{A}_{(ij)}^{-1} (h_{ij}^{(\delta)} - 1) \mathcal{L}_{(k)}^{-1} h_k^{(\delta)} \{h_k^{(\delta)-1}, V\} \right) + O((b^2 + c^2)^{3/2} \delta) \\
& = \left[-16\pi(8\pi\gamma^3 G\delta^3)^{-1} \sum_{ijk} \epsilon^{ijk} \text{Tr} \left(h_{ij}^{(\delta)} h_k^{(\delta)} \{h_k^{(\delta)}, V\} \right) \right] + O((b^2 + c^2)^{3/2} \delta). \tag{44}
\end{aligned}$$

Next, using $\Omega = -\sin \vartheta d\vartheta \wedge d\varphi \tau_3$, the Ω term in (40) becomes:

$$\int d^3x \epsilon_{ijk} \Omega_{ab}^i e^{-1} E^{aj} E^{bk} = -2 \int d^3x \text{Tr} (\Omega_{ab}^i \tau_i \epsilon_{jkl} e^{-1} E^{aj} E^{bk} \tau^l) \tag{45}$$

$$= -8(8\pi\gamma G\delta)^{-1} \int d^3x \sin \vartheta \text{Tr} (\tau_3 \mathcal{L}_x^{-1} h_x^{(\delta)} \{h_x^{(\delta)-1}, V\}) \tag{46}$$

$$= -32\pi(8\pi\gamma G\delta)^{-1} \text{Tr} (\tau_3 h_x^{(\delta)} \{h_x^{(\delta)-1}, V\}). \tag{47}$$

Set

$$C^{(\delta)} = \left[-2(\gamma^3 G\delta^3)^{-1} \sum_{ijk} \epsilon^{ijk} \text{Tr} \left(h_{ij}^{(\delta)} h_k^{(\delta)} \{h_k^{(\delta)}, V\} \right) \right] - \left[-4(\gamma G\delta)^{-1} \text{Tr} (\tau_3 h_x^{(\delta)} \{h_x^{(\delta)-1}, V\}) \right] \tag{48}$$

Then the classical Hamiltonian constraint is given by

$$\mathcal{C}_{\text{Ham}} = \lim_{\delta \rightarrow 0} C^{(\delta)} \tag{49}$$

Since all terms in $C^{(\delta)}$ are expressed purely in terms of our elementary variables—holonomies and triads—which have direct operator analogs, passage to quantum theory is now straightforward. We obtain:

$$\begin{aligned}
\hat{C}^{(\delta)} & = 2i(\gamma^3 \delta^3 \ell_{\text{Pl}}^2)^{-1} \text{Tr} \left(\sum_{ijk} \epsilon^{ijk} \hat{h}_i^{(\delta)} \hat{h}_j^{(\delta)} \hat{h}_i^{(\delta)-1} \hat{h}_j^{(\delta)-1} \hat{h}_k^{(\delta)} [\hat{h}_k^{(\delta)-1}, \hat{V}] + 2\gamma^2 \delta^2 \tau_3 \hat{h}_x^{(\delta)} [\hat{h}_x^{(\delta)-1}, \hat{V}] \right) \\
& = 4i(\gamma^3 \delta^3 \ell_{\text{Pl}}^2)^{-1} \left(8 \sin \frac{\delta b}{2} \cos \frac{\delta b}{2} \sin \frac{\delta c}{2} \cos \frac{\delta c}{2} \left(\sin \frac{\delta b}{2} \hat{V} \cos \frac{\delta b}{2} - \cos \frac{\delta b}{2} \hat{V} \sin \frac{\delta b}{2} \right) \right. \\
& \quad \left. + \left(4 \sin^2 \frac{\delta b}{2} \cos^2 \frac{\delta b}{2} + \gamma^2 \delta^2 \right) \left(\sin \frac{\delta c}{2} \hat{V} \cos \frac{\delta c}{2} - \cos \frac{\delta c}{2} \hat{V} \sin \frac{\delta c}{2} \right) \right) \tag{50}
\end{aligned}$$

(This is a special case of Eq. (26) in [30].)

The action of this operator on the eigenstates $|\mu, \tau\rangle$ of \hat{p}_b and \hat{p}_c is given by

$$\begin{aligned}
\hat{C}^{(\delta)} |\mu, \tau\rangle & = (2\gamma^3 \delta^3 \ell_{\text{Pl}}^2)^{-1} [2(V_{\mu+\delta, \tau} - V_{\mu-\delta, \tau}) \\
& \quad \times (|\mu + 2\delta, \tau + 2\delta\rangle - |\mu + 2\delta, \tau - 2\delta\rangle - |\mu - 2\delta, \tau + 2\delta\rangle + |\mu - 2\delta, \tau - 2\delta\rangle) \\
& \quad + (V_{\mu, \tau+\delta} - V_{\mu, \tau-\delta})(|\mu + 4\delta, \tau\rangle - 2(1 + 2\gamma^2 \delta^2)|\mu, \tau\rangle + |\mu - 4\delta, \tau\rangle)]. \tag{51}
\end{aligned}$$

However, while the classical constraint is a real function on the kinematical phase space Γ , \hat{C}^δ fails to be self-adjoint on the kinematical Hilbert space \mathcal{H} . We therefore need to add to it its Hermitian adjoint.⁹ The result is an operator $\hat{C}_{\text{grav}}^{(\delta)} := \frac{1}{2}(\hat{C}^{(\delta)} + \hat{C}^{(\delta)\dagger})$ given by:

$$\begin{aligned}
\hat{C}_{\text{grav}}^{(\delta)}|\mu, \tau\rangle &= (2\gamma^3\delta^3\ell_{\text{Pl}}^2)^{-1} [(V_{\mu+\delta, \tau} - V_{\mu-\delta, \tau} + V_{\mu+3\delta, \tau+2\delta} - V_{\mu+\delta, \tau+2\delta})|\mu + 2\delta, \tau + 2\delta\rangle \\
&\quad - (V_{\mu+\delta, \tau} - V_{\mu-\delta, \tau} + V_{\mu+3\delta, \tau-2\delta} - V_{\mu+\delta, \tau-2\delta})|\mu + 2\delta, \tau - 2\delta\rangle \\
&\quad - (V_{\mu+\delta, \tau} - V_{\mu-\delta, \tau} + V_{\mu-\delta, \tau+2\delta} - V_{\mu-3\delta, \tau+2\delta})|\mu - 2\delta, \tau + 2\delta\rangle \\
&\quad + (V_{\mu+\delta, \tau} - V_{\mu-\delta, \tau} + V_{\mu-\delta, \tau-2\delta} - V_{\mu-3\delta, \tau-2\delta})|\mu - 2\delta, \tau - 2\delta\rangle \\
&\quad + \frac{1}{2}(V_{\mu, \tau+\delta} - V_{\mu, \tau-\delta} + V_{\mu+4\delta, \tau+\delta} - V_{\mu+4\delta, \tau-\delta})|\mu + 4\delta, \tau\rangle \\
&\quad - (1 + 2\gamma^2\delta^2)(V_{\mu, \tau+\delta} - V_{\mu, \tau-\delta})|\mu, \tau\rangle \\
&\quad + \frac{1}{2}(V_{\mu, \tau+\delta} - V_{\mu, \tau-\delta} + V_{\mu-4\delta, \tau+\delta} - V_{\mu-4\delta, \tau-\delta})|\mu - 4\delta, \tau\rangle] . \tag{52}
\end{aligned}$$

Physical states in quantum theory are those which are symmetric under the ‘parity operator’ $\hat{\Pi}_b$ and lie in the kernel of the operator $\hat{C}_{\text{grav}}^{(\delta)}$. That is, the only non-trivial quantum Einstein’s equations are:

$$(\Psi|\hat{\Pi}_b = 0 \quad , \quad (\Psi|\hat{C}_{\text{grav}}^{(\delta)} = 0 \tag{53}$$

where $(\Psi|$ is an element of the dual Cyl_{KS}^* of the space Cyl_{KS} of finite linear combinations of almost periodic functions of b, c . Let us expand $(\Psi|$ using eigenbras $(\mu, \tau|$ (in Cyl_{KS}^*) of the triad operators:

$$(\Psi| = \sum_{\mu, \tau} \psi_{\mu, \tau} (\mu, \tau| \tag{54}$$

Then, $\psi_{\mu, \tau}$ can be regarded as wave functions in the triad (or, Riemannian geometry) representation. To exhibit the action of $\hat{C}_{\text{grav}}^{(\delta)}$ on $\psi_{\mu, \tau}$ we are led to separate the cases $\mu \geq 4\delta$ from $\mu < 4\delta$ because of the absolute-value $|\mu|$ in the volume eigenvalues. For $\mu \geq 4\delta$, quantum

⁹ Although the self-adjoint form of the Hamiltonian constraint was discussed briefly in [32], this point was ignored in the detailed treatment of [9]. A detailed treatment of the self-adjoint constraint and its semi-classical implications can be found in [26, 33]. It is used in a crucial way to obtain the physical Hilbert space in [10].

Einstein's equation (53) becomes

$$\begin{aligned}
\hat{C}_{\text{grav}}^{(\delta)} \psi_{\mu,\tau} &= 2\delta(\sqrt{|\tau+2\delta|} + \sqrt{|\tau|}) (\psi_{\mu+2\delta,\tau+2\delta} - \psi_{\mu-2\delta,\tau+2\delta}) \\
&\quad + (\sqrt{|\tau+\delta|} - \sqrt{|\tau-\delta|}) ((\mu+2\delta)\psi_{\mu+4\delta,\tau} - (1+2\gamma^2\delta^2)\mu\psi_{\mu,\tau} + (\mu-2\delta)\psi_{\mu-4\delta,\tau}) \\
&\quad + 2\delta(\sqrt{|\tau-2\delta|} + \sqrt{|\tau|}) (\psi_{\mu-2\delta,\tau-2\delta} - \psi_{\mu+2\delta,\tau-2\delta}) \\
&= 0.
\end{aligned} \tag{55}$$

For $\mu = 0$, we have

$$\sqrt{|\tau+2\delta|}\psi_{2\delta,\tau+2\delta} + (\sqrt{|\tau+\delta|} - \sqrt{|\tau-\delta|})\psi_{4\delta,\tau} - \sqrt{|\tau-2\delta|}\psi_{2\delta,\tau-2\delta} = 0; \tag{56}$$

for $\mu = \delta$:

$$\begin{aligned}
&2(\sqrt{|\tau+2\delta|} + \sqrt{|\tau|})\psi_{3\delta,\tau+2\delta} + 2(\sqrt{|\tau+2\delta|} - \sqrt{|\tau|})\psi_{\delta,\tau+2\delta} \\
&+ (\sqrt{|\tau+\delta|} - \sqrt{|\tau-\delta|}) (3\psi_{5\delta,\tau} + 2\psi_{3\delta,\tau} - (1+2\gamma^2\delta^2)\psi_{\delta,\tau}) \\
&- 2(\sqrt{|\tau-2\delta|} + \sqrt{|\tau|})\psi_{3\delta,\tau-2\delta} - 2(\sqrt{|\tau-2\delta|} - \sqrt{|\tau|})\psi_{\delta,\tau-2\delta} = 0,
\end{aligned} \tag{57}$$

for $\mu = 2\delta$:

$$\begin{aligned}
&(\sqrt{|\tau+2\delta|} + \sqrt{|\tau|})\psi_{4\delta,\tau+2\delta} - \sqrt{|\tau|}\psi_{0,\tau+2\delta} \\
&+ (\sqrt{|\tau+\delta|} - \sqrt{|\tau-\delta|}) (2\psi_{6\delta,\tau} + \psi_{2\delta,\tau} - (1+2\gamma^2\delta^2)\psi_{2\delta,\tau}) \\
&- (\sqrt{|\tau-2\delta|} + \sqrt{|\tau|})\psi_{4\delta,\tau-2\delta} + \sqrt{|\tau|}\psi_{0,\tau-2\delta} = 0,
\end{aligned} \tag{58}$$

and for $\mu = 3\delta$:

$$\begin{aligned}
&2(\sqrt{|\tau+2\delta|} + \sqrt{|\tau|})(\psi_{5\delta,\tau+2\delta} - \psi_{\delta,\tau+2\delta}) \\
&+ (\sqrt{|\tau+\delta|} - \sqrt{|\tau-\delta|}) (5\psi_{7\delta,\tau} + 2\psi_{\delta,\tau} - 3(1+2\gamma^2\delta^2)\psi_{3\delta,\tau}) \\
&- 2(\sqrt{|\tau-2\delta|} + \sqrt{|\tau|})(\psi_{5\delta,\tau-2\delta} - \psi_{\delta,\tau-2\delta}) = 0.
\end{aligned} \tag{59}$$

For values of μ not an integer multiple of δ one can derive similar expressions, but they will not be used in what follows. As remarked above, strictly speaking, the role of (55) is only to select physical wave functions. However, intuitively it can also be thought of as the quantum evolution equation. Recall that p_c can be regarded as an internal time in the classical theory of the model. Therefore, in the quantum theory, one may regard τ as internal time. Then, (55) can be interpreted as the quantum Einstein's equation, which evolves the wave functions in discrete steps of magnitude 2δ along τ .

So far, the parameter δ —and hence the size of the ‘time-step’ — is arbitrary. In the classical theory, we recover the Hamiltonian constraint only in the limit δ goes to zero. This is because, while the expression (4) is exact for all loops, curvature oF can be expressed in terms of holonomies only in the limit in which the areas of loops shrink to zero (see (43)). In quantum theory, however, the straightforward limit $\lim_{\delta \rightarrow 0} \hat{C}^{(\delta)}$ diverges (for the same reasons as in the isotropic case [9]). The viewpoint is that this occurs because the limit ignores the quantum nature of geometry, i.e., the fact that in full quantum geometry, the area operator has a minimum non-zero eigenvalue. The ‘correct’ quantization of the Hamiltonian constraint \mathcal{C}_{Ham} has to take in to account the quantum nature of geometry.

To do so, note first that the holonomies used in the expression of oF define quantum states with $\mu = \delta$ and $\tau = \delta$. Using these states, we can calculate the *quantum* area of each of the three faces of the elementary cell $\mathbb{S}^2 \times \mathcal{I}$, enclosed by the three curves —one along x , one along a longitude and one along the equator of \mathbb{S}^2 . The area operators defined by faces $S_{x,\vartheta}$, $S_{x,\varphi}$ and $S_{\vartheta,\varphi}$ are, respectively, $2\pi|\hat{p}_b|$, $2\pi|\hat{p}_b|$ and $\pi|\hat{p}_c|$ (see (14)). The quantum geometry states defined by the three holonomies are eigenstates of these area operators. Furthermore, they have the same eigenvalue: $\pi\gamma\delta\ell_{\text{Pl}}^2$. Now, in the full theory, we know that the area operator has a minimum non-zero eigenvalue, $a_o = 2\sqrt{3}\pi\gamma\ell_{\text{Pl}}^2$. The viewpoint, as in [9], is that in the calculation of the field strength oF , it is *physically* inappropriate to try to use surfaces of arbitrarily small areas. The best one can do is to shrink the area of the loop till it attains this quantum minimum a_o . This implies that we should set $\delta = 2\sqrt{3}$. To summarize, by using input from the full theory, we conclude that the quantum Hamiltonian constraint is given by $\hat{C}^{(\delta_o)}$ with $\delta_o = 2\sqrt{3}$.

Remarks:

i) Note that the character of the difference equation changes depending on whether $\mu \geq 4\delta$ or $\mu < 4\delta$. In particular, for $\mu \geq 4\delta$, the knowledge of the wave function at ‘times’ $\tau + 2\delta$ and τ determines only the combination $\psi_{\mu-2\delta,\tau-2\delta} - \psi_{\mu+2\delta,\tau-2\delta}$ at ‘time’ $\tau - 2\delta$ via (55). On the other hand, for $\mu = 0$ the wave function at $\tau + 2\delta$ and τ determines $\psi_{2\delta,\tau+2\delta}$ completely via (56). We will return to these differences in section IV C.

ii) As discussed in [9, 31, 34] in detail, one can recover the Wheeler-DeWitt equation with a specific factor ordering by taking a systematic limit of quantum Einstein’s equation. In the present model, the difference equation reduces to the following differential equation

for the wave function $\Psi(p_b, p_c) = \psi_{2p_b/\gamma\ell_{\text{Pl}}^2, p_c/\gamma\ell_{\text{Pl}}^2}$ on the classical minisuperspace \mathcal{M} :

$$16\ell_{\text{Pl}}^4 \left(\sqrt{p_c} \frac{\partial^2 \Psi}{\partial p_b \partial p_c} + \frac{p_b}{4\sqrt{p_c}} \frac{\partial^2 \Psi}{\partial p_b^2} + \frac{1}{2\sqrt{p_c}} \frac{\partial \Psi}{\partial p_b} \right) - 4 \frac{p_b}{\sqrt{p_c}} \Psi = 0 \quad (60)$$

For $\mu \gg \delta, \tau \gg \delta$, solutions of the discrete equation (55) can be approximated by those of this Wheeler DeWitt equation in a precise sense. In essence, this Wheeler-DeWitt equation is obtained from (55) by ignoring quantum geometry effects, i.e., by an appropriate $\delta \rightarrow 0$ limit. For the approximation by a differential equation, higher derivatives of the wave function ψ must be sufficiently small. Thus, the limit $\delta \rightarrow 0$ for the equation (55) exists only under additional assumptions on the wave function. It does not exist for the *operator* $\hat{C}_{\text{grav}}^{(\delta)}$ by itself.

C. Absence of Singularity

Recall that the classical singularity occurs at $p_c = 0$. The Wheeler DeWitt equation (60) is manifestly singular there. One can multiply it by $\sqrt{p_c}$ and define a new ‘internal time’, $\bar{p}_c = \ln p_c$ and obtain a regular equation in variables p_b, \bar{p}_c . However, since $\bar{p}_c \rightarrow -\infty$ at the singularity, this regular equation does not let us evolve across the singularity. Thus, the overall situation is the same as in the isotropic model.

To obtain the Wheeler-DeWitt equation from the ‘fundamental’ difference equation (55), we had to ignore quantum geometry effects. Thus, the key question is whether this failure of the Wheeler-DeWitt equation is an artifact of the approximation used. Can quantum geometry effects make a qualitative difference as they do in the isotropic model? We will now argue that the answer is in the affirmative.

Let us analyze the ‘evolution’ given by the difference equation (55). Does this evolution stop at $\tau = 0$? As in the isotropic case [7, 32], we will use the quantum constraint (55) as a recurrence relation starting at positive τ and evolve toward smaller values. However, now a new twist arises because for generic values $\mu \geq 4\delta$, (55) determines only the difference $\psi_{\mu+2\delta, \tau-2\delta} - \psi_{\mu-2\delta, \tau-2\delta}$ as a function of the initial values of the wave function. Therefore, quantum Einstein’s equation has to be supplemented with an appropriate boundary condition. From (55) which holds for $\mu \geq 4\delta$, one might conclude that, even if we restrict ourselves to the ‘lattice’ $\mu = n\delta$, one would have to specify the wave function at $\mu = 0, \delta, 2\delta, 3\delta$ at each ‘time step’. However, as remarked at the end of section IV B, the form of the difference

equations for $\mu < 4\delta$ is different. An examination of (56) and (57) reveals that it is in fact sufficient to specify the wave function $\psi_{\mu,\tau}$ just at $\mu = 0$ and $\mu = \delta$ at each ‘time-step’. This is a mathematically viable choice of the boundary condition and heuristically it corresponds to providing data the ‘horizon’. It will turn out that the issue of the resolution of singularity is insensitive to the precise choice of the boundary condition. We will therefore postpone the discussion of physically appropriate choices until after the main discussion.

Let us then start at some finite positive τ and carry out a backward evolution using (55). In contrast to the Wheeler-DeWitt equation (60), the coefficients in the difference equation (55) are always regular. At first one might think that this regularity by itself would be sufficient to ensure that the evolution would be well-defined across the singularity. Note however that this need not be the case. For, the backward evolution continues only as long as the coefficient of the wave function at $\tau - 2\delta$ is non-zero. The key question therefore is whether this coefficient vanishes. The issue is subtle. For instance, the coefficients in the non-self-adjoint operator $(\hat{C}^\delta)^\dagger$ are also non-singular. However, if one uses it in place of $\hat{C}_{\text{Ham}}^\delta$, one finds that the coefficient vanishes and one can not evolve across the singularity.

For our quantum Einstein’s equation (55), this coefficient is given by $\sqrt{|\tau - 2\delta|} + \sqrt{|\tau|}$. By inspection, *it never vanishes*. Therefore, the backward quantum evolution remains well-defined and determines the wave function not only for $\tau > 0$ but also in the new region with $\tau \leq 0$. In this precise sense, the classical black hole singularity can be traversed using quantum evolution and thus ceases to be a boundary of space-time.

Next, recall from section IV A that we could have begun with the expression (38) of the Hamiltonian constraint in the classical theory and then proceeded with quantization. What would be the status of the singularity resolution with this choice? The procedure to construct the quantum Hamiltonian constraint would be identical. However, the final result would be slightly different: γ^2 in (55) would be replaced by $1 + \gamma^2$. That is, in the new quantum Einstein equation, the coefficient of $\psi_{\mu,\tau}$ in (55) would be modified and the equation would become:

$$\begin{aligned}
& 2\delta(\sqrt{|\tau + 2\delta|} + \sqrt{|\tau|})(\psi_{\mu+2\delta,\tau+2\delta} - \psi_{\mu-2\delta,\tau+2\delta}) \\
& + (\sqrt{|\tau + \delta|} - \sqrt{|\tau - \delta|}) ((\mu + 2\delta)\psi_{\mu+4\delta,\tau} - (1 + 2(\gamma^2 + 1)\delta^2)\mu\psi_{\mu,\tau} + (\mu - 2\delta)\psi_{\mu-4\delta,\tau}) \\
& + 2\delta(\sqrt{|\tau - 2\delta|} + \sqrt{|\tau|})(\psi_{\mu-2\delta,\tau-2\delta} - \psi_{\mu+2\delta,\tau-2\delta}) \\
& = 0
\end{aligned} \tag{61}$$

(The same replacement holds for $\mu < 4\delta$.) Since the coefficient of $\psi_{\mu-2\delta,\tau-2\delta} - \psi_{\mu+2\delta,\tau-2\delta}$ is unaltered, one would again be able to ‘evolve’ across the singularity. Thus, the conclusion is robust within quantum geometry.¹⁰

To conclude, we will briefly return to the question of the boundary condition that are needed to make the evolution well-defined even away from the singularity. As noted earlier, in the backward evolution given by (56), $\psi_{2\delta,\tau_o}$ is determined by values of the wave function at $\tau > \tau_o$, so long as for $\tau_o \neq 0$. Furthermore, (55) then determines the wave function at $\mu = (2 + 4n)\delta$ for all $n \in \mathbb{N}$ and at that τ_o . One could now use pre-classicality arguments [35] (i.e. require that the wave function not oscillate on small scales) at large μ to choose the boundary values, eliminating entirely the need for specifying boundary conditions. While this could be reasonable in semiclassical regimes, for small $|\tau|$ it would be questionable to use pre-classicality in the μ -direction even if a pre-classical solution exist (which is not guaranteed, see e.g. the analysis in [36, 37]). More importantly, the slice $\tau = 0$ is special because $\psi_{2\delta,0}$ is not determined through (56) or otherwise. At this point, the condition (56) evaluated at $\tau = -2\delta$ gives instead a condition for previous values of the wave function, translating into one condition for the initial values. At $\tau = 0$, however, $\psi_{2\delta,0}$ has to be specified as boundary value in addition to $\psi_{0,0}$ because, unlike similar situations in isotropic models, it does not drop out of the evolution: it is needed in the recurrence (55) for $\mu = 4\delta$, $\tau = 2\delta$. There is thus no reduction in conditions on the wave function from the fact that $\psi_{2\delta,0}$ drops out of (56); a condition is simply transferred from boundary values to initial values. Note however that, even though the line $\tau = 0$ (corresponding to the classical singularity) does show special behavior, evolution does not break down there; it is only the boundary value problem which changes character.

While such a boundary value problem is mathematically well-defined, the resulting theory is not necessarily physically correct. For, in a physically interesting theory one would expect that, well away from the singularity, there should exist semi-classical solutions which are peaked on the classical trajectories. From Fig. 1 it is clear that the semi-classical solution peaked on the trajectory labelled by mass m would be sharply peaked at $\tau = \pm 4m^2$ on the line $\mu = 0$, whence semi-classical states peaked at different classical solutions will have

¹⁰ However, the relation between the wave function and the Wheeler–DeWitt approximation is less direct with this form of the constraint.

quite different forms near $\mu = 0$. Therefore, the theory obtained by simply fixing the wave function on (or near) $\mu = 0$ is not likely to admit a sufficiently rich semi-classical sector. Indeed, it is not clear that any boundary condition at (or near) $\mu = 0$ will be physically viable. Rather, one may have to impose the boundary condition at $\mu = \infty$, e.g., by requiring that the wave function should vanish there. Indeed, the form of dynamical trajectories of Fig. 1 implies that *every* semi-classical state would share this property. Moreover, the corresponding condition does hold in the closed isotropic model with a massless scalar field as source [38]. As in that case, one would expect that these semi-classical states would span a dense subspace in the physical Hilbert space. If so, requiring that the wave functions vanish at $\mu = \infty$ would be physically justified. Heuristics suggest that this strategy is viable for the Wheeler-DeWitt equation (60). However, detailed numerical analysis is necessary to establish that the difference equations (55) or (61) admits solutions satisfying this condition (or a suitable modification thereof) for all τ and that its imposition makes the evolution unambiguous.

V. DISCUSSION

Results of the last two sections support a general scenario that has emerged from the analysis of singularities in quantum cosmology. It suggests that the classical singularity does not represent a final frontier; the *physical* space-time does not end there. In the Planck regime, quantum fluctuations do indeed become so strong that the classical description breaks down. The space-time continuum of classical general relativity is replaced by discrete quantum geometry which remains regular during the transition through what was a classical singularity. Certain similarities between the Kantowski-Sachs model analyzed here and a cosmological model which has been studied in detail [10] suggest that there would be a quantum bounce to another large classical region. If this is borne out by detailed numerical calculations, one would conclude that quantum geometry in the Planck regime serves as a bridge between two large classical regions. Space-time may be much larger than general relativity has had us believe.

However, as indicated at the end of section IV B significant numerical work is still needed before one can be certain that this scenario is really borne out in the model. Moreover, this is a highly simplified model. It is important to check if the qualitative conclusions remain

robust as more and more realistic features are introduced. First, one should extend the analysis so that the space-time region outside the horizon is also covered. A second and much more important challenge is incorporation of an infinite number of degrees of freedom by coupling the model, e.g., to a spherically symmetric scalar field. First steps along these lines have been taken [39, 40] and one can see that the evolution still extends beyond the classical singularity [41]. But a comprehensive treatment still remains a distant goal.

Acknowledgements:

This work was supported in part by NSF grants PHY-0090091, and PHY-0354932, the Alexander von Humboldt Foundation, the C.V. Raman Chair of the Indian Academy of Sciences and the Eberly research funds of Penn State.

-
- [1] Misner C W 1972 Minisuperspaces, in *Magic without magic, John A. Wheeler festschrift* ed Klauder J R (Freeman, San Fransisco)
 - [2] Hartle J B 1995 Spacetime quantum mechanics and the quantum mechanics of spacetime, in *Gravitation and quantizations: Proceedings of the 1992 Les Houches summer school* (North Holland, Amsterdam) *Preprint gr-qc/9304006*
 - [3] Hartle J B and Hawking S W 1983 The wave function of the universe *Phys. Rev.* **D28** 2960-2975
 - [4] Ashtekar A and Lewandowski L 2004 Background independent quantum gravity: A status report *Class. Quant. Grav.* **21** R53-R152
 - [5] Rovelli C 2004 *Quantum Ggravity* (Cambridge University Press, Cambridge)
 - [6] Thiemann T 2004 *Modern canonical quantum general relativity* (Cambridge University Press, Cambridge, at press) *Preprint gr-qc/0110034*
 - [7] Bojowald M 2001 Absence of singularity in loop quantum cosmology *Phys. Rev. Lett.* **86** 5227–5230
 - [8] Bojowald M 2003 Homogeneous loop quantum cosmology *Class. Quantum Grav.* **20** 2595–2615
 - [9] Ashtekar A, Bojowald M and Lewandowski J 2003 *Mathematical structure of loop quantum cosmology* *Adv. Theor. Math. Phys.* **7** 233–268

- [10] Ashtekar A, Pawłowski T and Singh P 2005 Quantum nature of the big bang: An analytical and numerical study *Preprint* in preparation
- [11] Dafermos M 2003 Black hole formation from a complete regular past *Preprint* [gr-qc/0310040](#)
- [12] Horowitz G T and Myers R 1995 The value of singularities *Gen. Rel. Grav.* **27** 915-919
- [13] Marolf D 1995 Refined algebraic quantization: Systems with a single constraint *Preprint* [gr-qc/9508015](#); Giulini N and Marolf D 1999 On the Generality of refined algebraic quantization *Class. Quant. Grav.* **16** 2479–2488; A uniqueness theorem for constraint quantization *Class. Quant. Grav.* **16** 2489–2505
- [14] Ashtekar A and Bojowald M 2005 Black hole evaporation: A paradigm *Class. Quant. Grav.* **22** 3349–3362
- [15] Stevens C R, 't Hooft G and Whiting B F 1994 Black hole evaporation without information loss, *Class. Quantum. Grav.* **11** 621–648
- [16] Hayward S 2005 The disinformation problem for black holes (conference version), [gr-qc/0504038](#)
- [17] Modesto L 2004 The Kantowski-Sachs space-time in loop quantum gravity *Preprint* [gr-qc/0411032](#)
- [18] Kobayashi S and Nomizu K 1963 *Foundations of differential geometry* (John Wiley, New York)
- [19] Brodbeck O 1996 On symmetric gauge fields for arbitrary gauge and symmetry groups *Helv. Phys. Acta* **69** 321–324
- [20] Bojowald M and Kastrup H A 2000 Symmetry reduction for quantized diffeomorphism invariant theories of connections *Class. Quantum Grav.* **17** 3009–3043
- [21] Ashtekar A and Samuel J 1991 Bianchi cosmologies: The role of spatial topology *Class. Quant. Grav.* **8** 2191–2215
- [22] Kastrup H A and Thiemann T 1994 Spherically symmetric gravity as a completely integrable system *Nucl. Phys. B* **425** 665–686
- [23] Kuchař K V 1994 Geometrodynamics of Schwarzschild black holes *Phys. Rev. D* **50** 3961–3981
- [24] Thiemann T 1996 Anomaly-free formulation of non-perturbative, four-dimensional Lorentzian quantum gravity *Phys. Lett.* **B380** 257–264; 1998 Quantum spin dynamics (QSD) *Class. Quant. Grav.* **15** 839–873; 1998 QSD V : Quantum gravity as the natural regulator of matter quantum field theories *Class. Quant. Grav.* **15** 1281–1314
- [25] Ashtekar A, Fairhurst S and Willis J 2003 Quantum gravity, shadow states, and quantum

- mechanics *Class. Quantum Grav.* **20** 1031–1062
- [26] Willis J 2004 On the low-energy ramifications and a mathematical extension of loop quantum gravity *Ph.D. Dissertation, Penn State*
- [27] Bojowald M 2001 Inverse scale factor in isotropic quantum geometry *Phys. Rev.* **D64** 084018
- [28] Bojowald M 2005 Degenerate Configurations, Singularities and the Non-Abelian Nature of Loop Quantum Gravity *Preprint gr-qc/0508118*
- [29] Brunnemann J and Thiemann T 2005 Unboundedness of Triad-Like Operators in Loop Quantum Gravity *Preprint gr-qc/0505033*
- [30] Bojowald M, Date G and Vandersloot K 2004 Homogeneous loop quantum cosmology: The role of the spin connection *Class. Quantum Grav.* **21** 1253–1278
- [31] Bojowald M and Date G 2004 Consistency conditions for fundamentally discrete theories *Class. Quantum Grav.* **21** (2004) 121–143.
- [32] Bojowald M 2002 Isotropic loop quantum cosmology *Class. Quantum Grav.* **19** 2717–2741.
- [33] Ashtekar A, Bojowald M, Willis J 2005 Modifications to Friedmann’s equation from quantum geometry *Preprint in preparation*
- [34] Bojowald M 2001 The semiclassical limit of loop quantum cosmology *Class. Quantum Grav.* **18** L109–L116
- [35] Bojowald M 2001 Dynamical initial conditions in quantum cosmology *Phys. Rev. Lett.* **87** 121301
- [36] Cartin D and Khanna G 2005 Absence of pre-classical solutions in Bianchi I loop quantum cosmology *Phys. Rev. Lett.* **94** 111302
- [37] Date G 2005 Pre-classical solutions of the vacuum Bianchi I loop quantum cosmology *Preprint gr-qc/0505030*
- [38] Ashtekar A, Pawłowski T and Singh P 2005 Closed cosmology and quantum geometry *Preprint in preparation*
- [39] Bojowald M 2004 Spherically symmetric quantum geometry: states and basic operators *Class. Quantum Grav.* **21** 3733–3753; Bojowald M and Swiderski R 2004 The volume operator in spherically symmetric quantum geometry *Class. Quantum Grav.* **21** 4881–4900; Bojowald M and Swiderski R 2005 Spherically symmetric quantum geometry: Hamiltonian constraint *Preprint in preparation*
- [40] Husain V and Winkler O 2004 Quantum resolution of black hole singularities *Preprint*

gr-qc/0410125

- [41] Bojowald M 2005 Non-singular black holes and degrees of freedom in quantum gravity *Phys. Rev. Lett.* **95** 061301