Sixth International Conference on Gravitation and Cosmology IOP Publishing
Journal of Physics: Conference Series **140** (2008) 012007 doi:10.1088/1742-6596/140/1/012007

Quantum Gravity and the Information Loss Problem

Madhavan Varadarajan
Raman Research Institute, Bangalore 560 080
E-mail: madhavan@rri.res.in

Abstract. We review the standard picture of black hole evaporation and the ensuing Information Loss Problem. Next, we describe an alternative view of the evaporation process due to Ashtekar and Bojowald (AB) which rests on the assumption of singularity resolution in quantum gravity. To endow the AB proposal with precision as well as to test it in a quantitative setting, we consider the Callen- Giddings- Harvey- Strominger (CGHS) toy model of 2-dimensional black holes. We propose a quantum framework in which to understand this model and show that the proposed framework implies that the classical spacetime manifold acquires a quantum extension, as in the AB paradigm. We show that asymptotic analysis coupled with the quantum framework suggests a unitary picture of black hole evaporation consistent with early (on the extended null infinity) time Hawking radiation.

1. Introduction

The standard picture of black hole evaporation due to Hawking [1] implies that a black hole, formed through collapse of matter in a pure quantum state, evolves to a mixed state comprised primarily of thermal Hawking radiation. This violation of unitarity is known as the Information Loss Problem. It is widely accepted that the Information Loss Problem will find its resolution in a theory of quantum gravity. One such proposed resolution, due to Ashtekar and Bojowald [2] rests on the expectation that a theory of quantum gravity resolves the singularities of the classical theory. The key idea of Ashtekar and Bojowald is that since (quantum) physics does not end at the classical singularity, there is a quantum extension of spacetime beyond the singularity, and consequently, an extension of future null infinity as well. This extension provides a structure whereby asymptotic observers who wait "long enough" see purity of the outgoing state restored as a result of correlations between early Hawking radiation and the part of the incoming state which propagates through the classically singular region and emerges in the quantum extension.

Since no satisfactory theory of quantum gravity is available yet, the AB proposal is necessarily qualitative. In order to test their ideas as well as to sharpen their proposal, it is useful to examine the puzzle in the context of toy models which admit black hole solutions and in which progress can be made towards a complete quantization. Such a model is provided by the Callen- Giddings-Harvey- Strominger (CGHS) model [3] of 2-dimensional black holes. As we shall see, progress towards the non- perturbative quantization of this model is possible. The ensuing quantum framework, while by no means complete, is rich enough to provide an AB type resolution (with several attendant subtelities and insights of a robust enough nature that they apply to the realistic 4-dimensional case) to the information loss problem [4, 5].

We review the information loss problem in section 2 and the AB proposal in section 3. Sections 4 to 7 are devoted to a report on collaborative work with Ashtekar and Taveras [4, 5] on the

CGHS model. Section 8 contains our conclusions. Since a short summary of our work with essential technical details is already available [4] and a long, detailed account is in preparation [5], here we present an overview of the main results without proofs and technicalities. We use obvious notation wherever possible.

2. The Information Loss Puzzle

By definition, a classical black hole cannot radiate to future null infinity. However, in his seminal work [1], Hawking showed that black holes do radiate when leading order *quantum* effects are taken into account. Specifically, Hawking analysed the behaviour of quantum matter fields propagating on the classical spacetime geometry of a spherically symmetric black hole and showed that such a black hole spontaneously radiates particles of the quantum matter field at late times. The radiation is thermal at the 'Hawking' temperature T_H where $kT_H \sim (\frac{m_P}{M})m_P c^2 \sim \frac{1}{M}$. Here k is the Boltzmann constant, m_P the Planck mass and M the black hole mass. Since energy conservation demands that the black hole mass must decrease due to the loss by radiation, the black hole is said to 'evaporate'. Indeed, the relation $T_H \sim \frac{1}{M}$ suggests the sequence: Hawking radiation $\rightarrow$ black hole mass loss $\rightarrow$ higher Hawking temperature $\rightarrow$ more radiation.

The time scale for evaporation can be estimated as $T_{\text{evap}} = (\frac{1}{M}\frac{dM}{dt})^{-1}$ with $\frac{dM}{dt}$ estimated through the Stefan- Boltzmann Law as $\sigma T_H^4 R_S^2$ where $R_S \sim \frac{G}{c^2}M$ is the Schwarzschild radius of the black hole. The time scale over which the black holes settles down after it is perturbed can be estimated as the time it takes for light to travel across the black hole so that $T_{\text{settling}} \sim R_S/c$. This yields $T_{\text{settling}}/T_{\text{evap}} \sim m_P^2/M^2$. Thus, as long as $M >> m_P$, evaporation is a quasistatic process and the geometry can be modelled as that of a 1 parameter family of black holes of decreasing mass as depicted in Fig. 1.

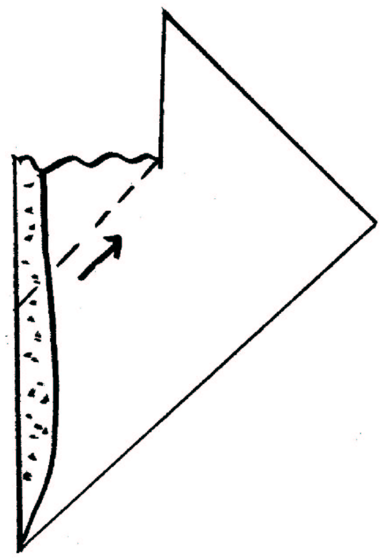

Figure 1. Figure caption

Evaporation continues till most of the mass (of the order of $M - m_P$) is lost as radiation, at which stage both the quasistatic approximation and the quantum field- on- curved- spacetime approximation break down. The collapsing matter admits an underlying quantum description and at early times this quantum state is pure. The final state after collapse is that of a 'remnant' of Planck size and lots of thermal radiation. The latter being a mixed state implies that the final state is pure only if the remnant has enough correlations with the radiation; this is ruled out by low energy physics arguments [6] and we have the puzzle that information has been lost (i.e. quantum evolution is not unitary) during the evaporation process.

Despite its robustness, there is a loophole in the above argument: while the evaporation process is very slow, the black hole lifetime is very large and small corrections to quasistaticity can have a cumulative effect such that there is a significant departure from the above picture after enough time has elapsed. Hence there is room for alternative pictures of the evaporation process. We describe one such alternative which was proposed by Ashtekar and Bojowald in Ref. [2].

3. The Ashtekar Bojowald Paradigm

In Fig. 1, the final state is not just a remnant and radiation; there is also a (singular) boundary of spacetime. Ashtekar and Bojowald (AB) note that the notion of classical spacetime should break down not only at the end point of evaporation but all along the singularity. Using results from symmetry reduced models, they suggest that the classical singularity will be resolved also in the full theory of quantum gravity. The picture which AB put forward has the following features [2]:
(a) Classical spacetime is not a viable concept near the singularity and is replaced by some, as yet unknown, quantum construct.
(b) Quantum evolution (of gravity + matter) is well defined through the classical singularity. Thus, spacetime does not end at the classical singularity but admits a quantum extension.
(c) Quantum evolution is unitary. The 'missing' information is recovered from the correlations between the Hawking radiation and the quantum fields which re- emerge on the 'other side' of the classical singularity.
(d) Since the classical geometry near the horizon is not to be trusted, and since the event horizon is a global construct which is sensitive to the entire spacetime geometry (including near the singularity), [1] it follows that the event horizon may not be a robust construct when quantum effects are accounted for. Therefore AB characterise the evaparation process through a quasilocal construct called the Dynamical Horizon. A dynamical horizon [8] is a 3-dimensional hypersurface foliated by marginally trapped surfaces. It has been shown [8] that: (i) if (positive stress energy) matter falls into a dynamical horizon, the horizon becomes spacelike and its area increases, (ii) if there is no matter infall the horizon is null and of constant area, and (iii) if matter is emitted by a dynamical horizon, it becomes timelike and its area decreases.

From (a)- (d), AB are led to the spacetime diagram of Fig. 2. Infalling matter collapses and a (spacelike) dynamical horizon forms. Once all the collapsing matter has entered the dynamical horizon, the horizon becomes null and looks like (part of) a stationary event horizon. Hawking radiation starts and the dynamical horizon becomes timelike. What evaporates is the area of the dynamical horizon. The classical singularity is replaced by a quantum region and the classical spacetime obtains a quantum extension. Asymptotic observers who live long enough see the purity of the outgoing state being restored by virtue of correlations between the the part of the state which evolves through the classically singular region, and the Hawking radiation.

4. The CGHS model

4.1. Classical Theory

The CGHS action [3] depends on a 2-dimensional metric, g_{ab}, a dilaton field, ϕ, and a scalar field, f. In units wherein the speed of light is unity, the action is

$$S = \frac{1}{2G}\int d^2x\sqrt{g}e^{-2\phi}[R + 4(\nabla\phi)^2 + 4\kappa^2] - \frac{1}{2}\int d^2x\sqrt{g}g^{ab}\nabla_a f\nabla_b f \tag{1}$$

[1] Indeed, Hajicek [7] has shown that a modification of the geometry within a planck length of the singularity has the ability to make the event horizon disappear.

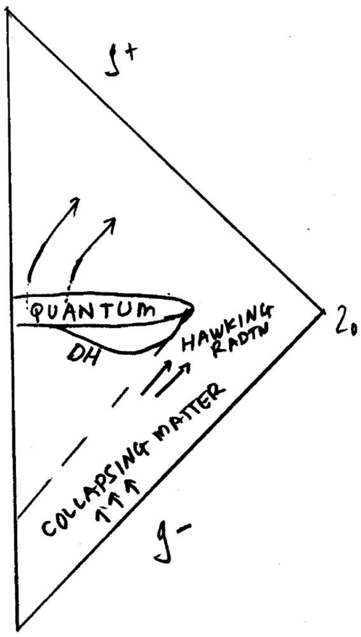

Figure 2. AB Picture

Here κ is a 'cosmological' constant of dimension L^{-1} and G is the 2-dimensional "Newton's" constant of dimension $L^{-1}M^{-1}$.

The model is classically exactly solvable. The general solution is as follows. Since spacetime is 2-dimensional, the metric is conformally flat so that $g_{ab}dz^a dz^b = -\frac{dz^+ dz^-}{\Omega}$, where $z^{\pm} = t \pm z$ and Ω^{-1} is the conformal factor. Since the scalar field is conformally coupled, it satisfies the flat pacetime wave equation $\partial_- \partial_+ f = 0$ so that f is of the form $f = f_{(+)}(z^+) + f_{(-)}(z^-)$. The functions $f_{(+)}$ and $f_{(-)}$ are called 'left' and 'right' movers. It turns out that $(\Omega - 1)$ is given by double integrals of the stress energy of f. The dilaton can also be expressed in terms of f. Hence $f_{(\pm)}(z^{\pm})$ parametrise the true degrees of freedom of the model.

Clearly , if $f = 0$, $\Omega = 1$ and we have flat spacetime. The solution corresponding to matter collapse is depicted in Fig. 3 and is obtained when $f_{(-)} = 0$ and $f_{(+)}$ is chosen to be some function of compact support The spacetime ends at the singularity as a result of which the spacetime manifold is a proper subset of the full (z^+, z^-) plane. The last light ray to reach right future null infinity traces out the event horizon. It may seem that there is an event horizon and a black hole with respect to left future null infinity as well. However, it turns out that whereas right future null infinity is geodesically complete, left future null infinity is not. Hence the singularity is left naked and there is a black hole structure only from the right (Geodesic completeness of null infinity is an essential ingredient for the presence of a black hole [9]).

For the passage to quantum theory it is useful to make a change of variables in the dilaton-metric sector from ϕ, Ω to the new variables Φ, Θ defined through:

$$\Phi := e^{-2\phi} \quad \Omega := \Theta^{-1}\Phi. \tag{2}$$

The classical equations of motion for the variables (Φ, Θ, f), in the $z^{\pm}$ coordinates are:

$$\begin{aligned} \partial_+ \partial_- f &= 0 \\ \partial_+ \partial_- \Phi + \kappa^2 \Theta &= G T_{+-} \\ \Phi \partial_+ \partial_- \ln \Theta &= -G T_{+-} \end{aligned} \tag{3}$$

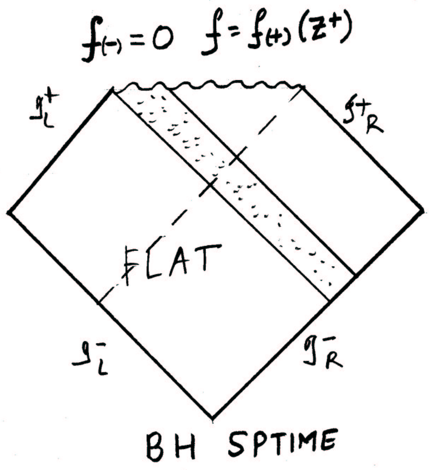

Figure 3. Classical solution

$$
\begin{aligned}
-\partial_+^2 \Phi + \partial_+ \Phi \partial_+ \ln \Theta &= GT_{++} \\
-\partial_-^2 \Phi + \partial_- \Phi \partial_- \ln \Theta &= GT_{--}
\end{aligned}
\tag{4}
$$

where T_{+-}, T_{++}, T_{--} are the $z^{\pm}$ components of the stress energy tensor of f. If (4) are imposed at past (right and left) null infinity, they are propagated by (3). Therefore we will refer to (3) as *dynamical equations* and ensure that (4) are satisfied by choosing appropriate boundary conditions at past null infinity. By virtue of the conformal coupling the classical stress energy tensor takes the standard flat spacetime form so that $T_{+-} = 0$ and $T_{\pm\pm}$ are quadratic in $\partial_{\pm} f$ and independent of the conformal factor Ω which determines the curved spacetime metric g_{ab}.

4.2. Hawking radiation

Since we only have a black hole with respect to right future null infinity, we shall be interested in the investigation of Hawking radiation near this segment of null infinity. Accordingly [10], it can be shown that a Hawking type analysis based on mode expansions of a (right moving, conformally coupled) quantum *test* field on the background spacetime of Fig. 3 implies that at late times the black hole emits thermal radiation to right future null infinity. Specifically, one chooses the quantum state of the test field to be its vacuum with respect to inertial observers near left past null infinity, traces over modes which fall into the singularity and investigates the particle content of the resulting density matrix with respect to inertial observers near right future null infinity to conclude that they observe thermal radiation at late times. In contrast to the case of 4d general relativistic collapse, the temperature of this thermal Hawking radiation is independent of its mass and proportional to κ. Nevertheless, the black hole does evaporate and there is an ensuing information loss problem.

The same conclusion may be reached via an analysis of the stress energy tensor operator of the test field. Since this operator involves the product of test field operators at the same spacetime point, it can only be defined after a regularization procedure. The standard point splitting regularization [11] brings the arguments of the test field operators to coincidence along a geodesic of the background spacetime and subtracts out an appropriate divergent piece. The resulting expectation value of the stress energy tensor is a sum of the classical stress energy and a quantum correction of order $\hbar$. For example (the expectation value of) the trace of the stress energy tensor does not vanish;it is equal (upto a numerical factor) to $\hbar$ times the scalar curvature of the background spacetime. Setting the background geometry to be that of the CGHS black hole, the stress energy tensor expectation value (in the vacuum state for the inertial observers near left past null infinity) can be computed at late times near right future null infinity. It turns out [3, 10] that the stress energy flux is thermal at the Hawking temperature consistent with the mode analysis calculation.

4.3. Quantum Framework.

Consider the operator versions of the classical equations. The matter field satisfies the flat spacetime wave equation $\partial_+\partial_-\hat{f} = 0$. Hence we quantize $\hat{f}(z^+, z^-)$ in the standard flat spacetime Fock representation. Thus $\hat{f}(z^+, z^-) = \hat{f}_{(+)}(z^+) + \hat{f}_{(-)}(z^-)$ with

$$\hat{f}_{(\pm)} = \int_0^\infty \frac{dk}{\sqrt{4\pi k}}[\hat{a}_{(\pm)}(k)e^{\mp ikz^\pm} + \hat{a}^\dagger_{(\pm)}(k)e^{\pm ikz^\pm}], \tag{5}$$

where $\hat{a}_{(\pm)}(k)$ are the annihilation operators for the right and left moving modes. The Fock space is then a product of right and left moving Fock spaces. This Fock space is the Hilbert space of the quantum theory. Although the quantum states in this Hilbert space seem to be those of the scalar field, it is important to remember that they are quantum states of the full gravity-dilaton- scalar field system; it is only that we have chosen to parametrise the true degrees of freedom by the scalar field. Thus the remaining operators of the theory, $\hat{\Theta}, \hat{\Phi}$ are to be defined on this Fock space.

As in the classical theory, we expect that $\hat{\Theta}, \hat{\Phi}$ will be determined by $\hat{f}(z^+, z^-)$ via operator versions of the classical equations. It is here that our quantization is incomplete: we do not yet know how to define the operator equations of motion. The missing element is the following. The classical equations involve the stress energy tensor and hence the quantum ones involve the stress energy operator. As already noted, in the context of quantum fields on classical spacetime, the classical spacetime plays a crucial role in the definition of the (test field) stress energy operator expectation value. In contrast, here we have a *quantum spacetime* characterised by the operator $\hat{\Omega}$ so that we need a definition of "quantum stress energy on *quantum* spacetime". While such a notion is yet undeveloped, progress seems possible based on developments in algebraic quantum field theory [12] and we intend to work on this issue in the future. Here, we shall *assume* that some satisfactory notion of the stress energy operator on a quantum spacetime can be developed in such a way that operator equations for $\hat{\Theta}, \hat{\Phi}$ can be defined and solved so that $\hat{\Theta}, \hat{\Phi}$ are determined by $\hat{f}_{(\pm)}$ as operators on the Hilbert space of the system. As we shall see, even this bare- bones framework enables us to make significant progress with respect to the issue of information loss, the reason being that, as evinced by equation (5), the natural arena for quantum theory is now the *entire* (z^+, z^-) plane.

The next step is to fix a "quantum black hole state" in the Fock space. Since classical black holes are formed by matter data $f_- = 0$, so that $f = f_+(z^+)$, we choose the quantum state to be a coherent state (with respect to the mode operators $\hat{a}_{(\pm)}(k)$) patterned on this data so that the state is a product of the '-' vacuum state with the '+' coherent state based on $f_{(+)}$. We denote this state by $|0_- > |f_{(+)} >$. By virtue of the fact that there is black hole behaviour only with respect to right future null infinity, the Information Loss Problem takes the form: *What happens to the* $|0_- >$ *part of the quantum black hole state during black hole evaporation?.*

5. Trial solution to the operator equations

To develop intuition about the operator equations, we use the only metric available, namely the auxiliary flat metric dz^+dz^-. The stress energy tensor defined with respect to this metric is the standard normal ordered (with respect to $\hat{a}_{(\pm)}(k), \hat{a}^\dagger_{(\pm)}(k)$) one (with vanishing trace) and the operator equations can be exactly solved for the operators $\hat{\Theta}, \hat{\Phi}$ in terms of $\hat{f}$. It turns out that the operators $\hat{\Theta}, \hat{\Phi}, \hat{\Omega}(z^+, z^-) := \hat{\Phi}\hat{\Theta}^{-1}$ are defined everywhere on the (z^+, z^-) plane with $\hat{\Theta}$ equal to a c-number function times identity. The expectation value of $\hat{\Omega}$ in the black hole state equals the classical conformal factor appropriate to the data $(f_{(-)} = 0, f = f_{(+)})$ below the singularity, vanishes on the singularity and is well defined "above the singularity" (see Fig. 4) so

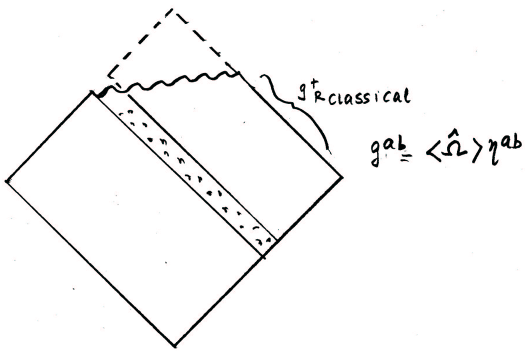

Figure 4. Expectation value geometry

that there is a quantum extension of the classical spacetime manifold to the full (z^+, z^-) plane. It can be shown [5] that fluctuations of (suitably smeared) $\hat{\Omega}$ are very large near the classical singularity so that the classical "expectation value" geometry is not a physically appropriate description there.

Next, the particle content of the '-' part of the black hole quantum state as viewed by asymptotic observers of the expectation value (classical) geometry near the classical future right null infinity can be examined. In other words one asks for the particle content of the state $|0_- >$ as observed by asymptotic inertial observers at right future null infinity of the classical geometry. It can be shown that such observers at late times see exactly a Hawking flux of particles at the mass- independent Hawking temperature [5]! Note that $|0_- >$ is *not* a test field; it is the right moving part of the full non-perturbative quantum state of the matter- dilaton- gravity system.

While the classical geometry and Hawking radiation can be recovered in this trial solution, there is no back reaction of the radiation on the geometry by virtue of the physically inappropriate choice of the auxiliary flat metric to define the stress energy operator. Nevertheless, the trial solution already exhibits key aspects of the AB paradigm such as the inappropriateness of the classical description along the singularity as well as a quantum extension of spacetime. To reiterate: Whereas the classical (inverse) conformal factor (equal to $\langle\hat{\Omega}\rangle$) vanishes at the singularity, the corresponding quantum operator $\hat{\Omega}$ is still well defined there. There is no physical problem at the classical singularity; it is simply a place where the classical description (by $\langle\hat{\Omega}\rangle$) breaks down. Due to the large fluctuations of $\hat{\Omega}$ in the classically singular region, the expectation value $\langle\hat{\Omega}\rangle$ does not capture the true physics. There is a quantum extension of the classical spacetime, beyond the classical singularity, to the entire $z^+ - z^-$ plane and the quantum metric, $\hat{g}^{ab} = -2\hat{\Omega}((\frac{\partial}{\partial z^+})^a(\frac{\partial}{\partial z^-})^b + (\frac{\partial}{\partial z^+})^b(\frac{\partial}{\partial z^-})^a)$, is a globally (i.e. on the entire $z^+ - z^-$ plane) well defined solution of the operator equations.

6. Mean Field Approximation (MFA).

Since a modification of the classical geometry by back reaction is a key feature of the evaporation process, we study the operator equations in an approximation which accommodates back reaction. The approximation (which we call the MFA) is one in which we neglect fluctuations of $\hat{\Theta}, \hat{\Phi}$ but not $\hat{f}$ i.e. the dilaton- geometry fluctuations are neglected whereas the matter fluctuations are retained.[2] (In full blown 4d quantum gravity such an approximation leads to the much studied semiclassical Einstein equations, $G_{ab} = 8\pi G < T_{ab} >$ which contain backreaction).

The resulting equations have been studied both analytically and numerically [13] with the following caveat. In contrast to our choice of smooth $f_{(+)}$, in Ref. [13] $f_{(+)}$ is chosen to be a shock wave with singular support along its line of infall. We shall assume here that the broad features of the MFA collapse of Ref. [13] continue to hold for smooth $f_{(+)}$. We describe these features below and depict them in Fig 5.

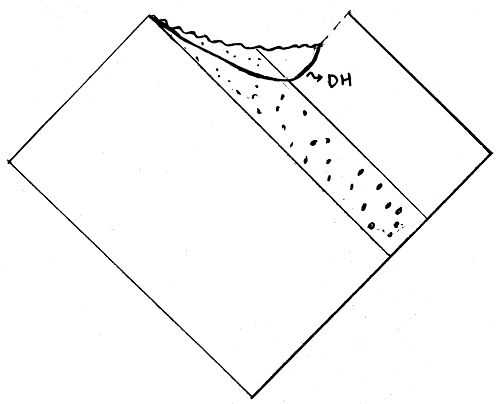

Figure 5. MFA Picture

The spacetime is still singular. A dynamical horizon forms which is initially spacelike and then turns around and becomes timelike. A Cauchy horizon develops in the causal shadow of the singularity beyond which evolution is not possible. Indications are that almost the entire mass of the classical spacetime is radiated away through Hawking radiation. Again, aspects of the AB paradigm, namely the appearance of a dynamical horizon are manifest. [3]

7. Asymptotic Analysis.

Our viewpoint towards the MFA singularity is that it is a place where the MFA breaks down; intuition from the trial solution suggests that the full quantum equations are still well defined there and that the singular region is replaced by one of quantum spacetime. While we expect the MFA to break down near regions of large curvature, we expect that the MFA should still be valid at right future null infinity of the physical (as opposed to auxiliary flat) spacetime. Hence, we perform an asymptotic analysis of the MFA equations near right future null infinity of the physical spacetime [4]. Note that since the classical description is expected to be good at early times, this null infinity coincides initially with that of the classical (and auxiliary flat) spacetime. It turns out that the knowledge of the non- perturbative quantum black hole state together with the MFA equations near the right future null infinity of the physical spacetime strongly constrain the response of the asymptotic geometry to the energy flux there.

[2] Such an approximation can be justified in the context of a large number, N, of conformally coupled scalar fields [3] in contrast to the $N = 1$ case studied above. However our framework can easily be extended to the case of N scalar fields the only change being that we have N copies of the Fock space with $N - 1$ of the fields chosen in their vacuum states and the remaining one in the quantum black hole state defined above.

[3] Since the spacetime is 2-dimensional, it is not possible to define trapped surfaces in the usual way; instead the dynamical horizon is defined through the behavior of the dilaton [13] drawing upon the analogy with a recasting of spherically symmetric 4d gravity as a 1+1 theory.

Journal of Physics: Conference Series **140** (2008) 012007 doi:10.1088/1742-6596/140/1/012007

Specifically, if we assume that the Hawking flux switches off smoothly at some finite affine parameter value along physical right future null infinity, then it follows *essentially uniquely* that the physical right future null infinity *coincides* with the that of the auxiliary flat spacetime! This is additional evidence of a quantum extension of the classical spacetime. (It seems physically reasonable to expect that the Hawking radiation switches off or decreases in such a way that only a finite amount of mass is radiated away- here we further assume that the switch off is smooth and that it occurs at "finite time"). Moreover, one finds (under this assumption) that $|0_- >$ is a *normalised, pure* state in the Hilbert space of inertial observers near right future null infinity of this (extended) physical spacetime. These observers see the state populated by an even number of particles; nevertheless the purity and unit norm of the state ensures that there is no information loss.

8. The Final Picture

Putting together all the ingredients discussed above, the spacetime diagram of the evaporating CGHS black hole is depicted in Fig. 6. The region interior to the past of the MFA singularity is

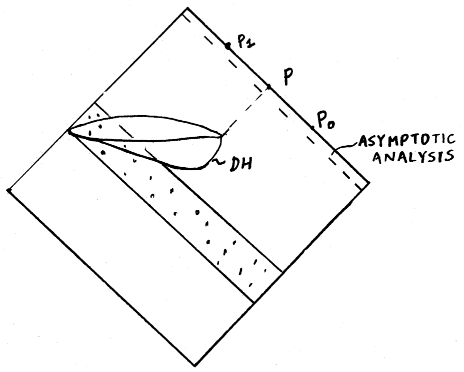

Figure 6. Final Picture

given by the MFA solution and that near future null infinity by asymptotic analysis. Intuition from the trial solution as well as the conceptual underpinnings of our quantum framework lead us to expect that the region around the MFA singularity is replaced by a quantum spacetime region (depicted by the oval region above the dynamical horizon) and that the spacetime is extended beyond this region so that it covers the entire (z^+, z^-) plane. By assumption the Hawking flux vanishes in the far future on right future null infinity of the spacetime.

The Information Loss Problem (see the last line of section 4) is resolved as follows (see Fig. 6). As mentioned above $|0_- >$ is a pure state in the Hilbert space of inertial observers near future right null infinity There is no significant energy flux beyond the point P of Fig. 6 (this is suggested by the MFA numerics [13] in conjunction with the assumption that the radiated energy does not exceed the total mass of the black hole;) and no remnant with a large number of internal states. Despite the absence of energy flux in the future, there are still the standard flat spacetime vacuum correlations between scalar field operators at early and late times on future null infinity i.e. $< 0_-|\partial_-\hat{f}_{(-)}(P_1), \partial_-\hat{f}_{(-)}(P_0)|0_- > \neq 0$ (see Fig. 6).

While the analysis of information loss in the language of correlations shows that the correlations continue to emerge at late times despite the absence of energy flux, most discussions of information loss are phrased in terms of entropy (as opposed to correlations). The definition of entropy requires us to "trace over particle modes" beyond some point on future null infinity. Since quantum fields *cannot be localised*, any definition of such modes necessarily involves a spread in their position. It is possible to localise modes with a spread by considering wave packets peaked at values of position and momenta. Choose some point P' on future right null infinity and trace over modes to the future of P'. This yields a density matrix whose entropy $S(P')$ can be estimated. It can be argued [5] that the entropy starts out at zero when P' is in the remote past, then starts increasing and then decreases. The decrease is due to correlations between particles emitted at early and later times. By the time $P' = P$ the entropy decreases close to zero. This is consistent with the fact that correlations between field operators at P_1 and P_0 never vanish, precisely due to the presence of spreads in the definition of modes.

In conclusion, aspects of a putative non-perturbative quantization, in conjunction with MFA numerics and asymptotic analysis point to a unitary picture of black hole evaporation with the following key features which seem robust enough to carry over to higher dimensions: (i) Singularity resolution (this is from intuition from the trial solution and is as yet, only an informed guess as to what will happen) (ii) Extension of classical spacetime (this is due to the quantum arena being the entire (z^+, z^-) plane and is further supported by asymptotic analysis) (iii) The inadmissibility of the notion of classical empty spacetime: there is always an underlying quantum state which supports field correlations even in the absence of stress energy.

References

[1] S. W. Hawking, *Comm. Math. Phys.***43**, 199 (1975).
[2] A. Ashtekar and M. Bojowald, *Class.Quant.Grav.***22**, 3349 (2005).
[3] C. G. Callan, S. B. Giddings, J. A. Harvey and A. Strominger, *Phys.Rev.***D45**, 1005 (1992).
[4] A. Ashtekar, V. Taveras and M. Varadarajan, e-Print:gr-qc 0801.1811
[5] M. Varadarajan, A. Ashtekar and V. Taveras (in preparation)
[6] J. Preskill, in *Proceedings of the International Symposium on Black Holes, Membranes, Wormholes, and Superstrings, Woodlands, TX, 1992*, (World Scientific, Houston, TX, 1992).
[7] P. Hajicek, *Phys.Rev.***D36**, 1065 (1987).
[8] A. Ashtekar and B. Krishnan, *Phys.Rev.***D68**, 104030 (2003).
[9] R. Geroch and G. T. Horowitz, *Phys. Rev. Lett.* **40**, 203 (1978).
[10] S. B. Giddings and W. M. Nelson, *Phys.Rev.***D46**, 2486 (1992).
[11] P.C.W. Davies, S.A. Fulling and W.G. Unruh, *Phys.Rev.***D13** 2720 (1976).
[12] M. Dütsch and K. Fredenhagen, *Commun. Math.Phys.* **203**, 71 (1999).
[13] L. Susskind and L. Thorlacius, *Nucl.Phys.***B382** 123 (1992); D. Lowe, *Phys. Rev.***47**, 2446 (1992); T. Priran and A. Strominger, *Phys. Rev.***48**, 4729 (1993).