Gravitational amplitudes in black hole evaporation: the effect of non-commutative geometry

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Received 21 July 2006, in final form 8 September 2006
Published 11 October 2006
Online at stacks.iop.org/CQG/23/6425

Abstract
Recent work in the literature has studied the quantum-mechanical decay of a Schwarzschild-like black hole, formed by gravitational collapse, into almost-flat spacetime and weak radiation at a very late time. The relevant quantum amplitudes have been evaluated for bosonic and fermionic fields, showing that no information is lost in collapse to a black hole. On the other hand, recent developments in non-commutative geometry have shown that, in general relativity, the effects of non-commutativity can be taken into account by keeping the standard form of the Einstein tensor on the left-hand side of the field equations and introducing a modified energy–momentum tensor as a source on the right-hand side. The present paper, relying on the recently obtained non-commutativity effect on a static, spherically symmetric metric, considers from a new perspective the quantum amplitudes in black hole evaporation. The general relativity analysis of spin-2 amplitudes is shown to be modified by a multiplicative factor \( F \) depending on a constant non-commutativity parameter and on the upper limit \( R \) of the radial coordinate. Limiting forms of \( F \) are derived which are compatible with the adiabatic approximation here exploited. Approximate formulae for the particle emission rate are also obtained within this framework.

PACS numbers: 03.70.+k, 04.60.Ds, 04.70.Dy

1. Introduction
Theoretical research in black hole physics has witnessed, over the last few years, an impressive amount of new ideas and results on at least four main areas.
(i) The problem of information loss in black holes, after the suggestion in [1] that quantum gravity is unitary and information is preserved in black hole formation and evaporation.

(ii) The related series of papers in [2–10], concerned with evaluating quantum amplitudes for transitions from initial to final states, in agreement with a picture where information is not lost, and the end state of black hole evaporation is a combination of outgoing radiation states.

(iii) The approach in [11–13], according to which black holes create instead a vacuum matter charge to protect themselves from the quantum evaporation.

(iv) The work in [14] where the authors, relying upon the previous findings in [15], consider a non-commutative radiating Schwarzschild black hole, and find that non-commutativity cures usual problems encountered in trying to describe the latest stage of black hole evaporation.

We have been therefore led to study how non-commutativity would affect the analysis of quantum amplitudes in black hole evaporation performed in [2–10]. Following [14], we assume that non-commutativity of spacetime can be encoded in the commutator of operators corresponding to spacetime coordinates, i.e. (the integer \( D \) below is even)

\[
[x^\mu, x^\nu] = i \theta^{\mu\nu}, \quad \mu, \nu = 1, 2, \ldots, D,
\]

where the antisymmetric matrix \( \theta^{\mu\nu} \) is taken to have a block-diagonal form

\[
\theta^{\mu\nu} = \text{diag}(\theta_1, \ldots, \theta_{D/2}),
\]

with

\[
\theta_i = \theta \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \quad \forall i = 1, 2, \ldots, D/2,
\]

the parameter \( \theta \) having dimension of length squared and being constant. As shown in [15], the constancy of \( \theta \) leads to a consistent treatment of Lorentz invariance and unitarity. The authors of [14] solve the Einstein equations with mass density of a static, spherically symmetric, smeared particle-like gravitational source as (hereafter we work in \( G = c = \hbar = 1 \) units)

\[
\rho_\theta(r) = \frac{M}{(4\pi \theta)^2} e^{-\frac{r^2}{4\theta}},
\]

which therefore plays the role of matter source. Their resulting spherically symmetric metric is

\[
ds^2 = -\left[ 1 - \frac{4M}{r} \gamma \left( \frac{3}{2}, \frac{r^2}{4\theta} \right) \right] \, dt^2 + \left[ 1 - \frac{4M}{r} \gamma \left( \frac{3}{2}, \frac{r^2}{4\theta} \right) \right]^{-1} \, dr^2 + r^2 (d\theta^2 + \sin^2 \theta \, d\phi^2),
\]

where we use the lower incomplete gamma function [14]

\[
\gamma \left( \frac{3}{2}, \frac{r^2}{4\theta} \right) = \int_0^{r^2} \sqrt{t} \, e^{-t} \, dt.
\]

In this picture, we deal with a mass distribution

\[
m(r) = \frac{2M}{\sqrt{\pi}} \gamma \left( \frac{3}{2}, \frac{r^2}{4\theta} \right),
\]

while \( M \) is the total mass of the source [14]. This mass function satisfies the equation

\[
m'(r) = 4\pi r^2 \rho_\theta(r),
\]

formally analogous to the general relativity case [10].
The work in [2–10] studies instead the quantum-mechanical decay of a Schwarzschild-like black hole, formed by gravitational collapse, into almost-flat spacetime and weak radiation at a very late time. The spin-2 gravitational perturbations split into parts with odd and even parity, and one can isolate suitable variables which can be taken as boundary data on a final spacelike hypersurface $\Sigma_{1F}$. The main idea is then to consider a complexified classical boundary-value problem where $T$ is rotated into the complex: $T \rightarrow |T|e^{-i\alpha}$, for $\alpha \in [0,\pi/2]$, and evaluate the corresponding classical Lorentzian action $S^{(2)\text{class}}_{\Sigma_{1F}}$ to quadratic order in metric perturbations. The genuinely Lorentzian quantum amplitude is recovered by taking the limit as $\alpha \rightarrow 0^+$ of the semiclassical amplitude $e^{iS^{(2)\text{class}}_{\Sigma_{1F}}}$ [2, 6, 10].

Section 2 studies the differential equations obeyed by radial modes within the framework of the adiabatic approximation, and section 3 obtains the resulting orthogonality relation in the presence of a non-vanishing non-commutativity parameter $\theta$. Section 4 derives the effect of $\theta$ on the expansion of the pure-gravity action functional, which can be used in the evaluation of quantum amplitudes along the lines of [2–10]. Absorption and emission spectra of a non-commutative Schwarzschild-like black hole are studied in section 5, while concluding remarks are presented in section 6, and relevant details are given in the appendix.

2. Equations for radial modes

The analysis in [10] holds for any spherically symmetric Lorentzian background metric

$$ds^2 = -e^{2(a(r,t))}dt^2 + e^{2(b(r,t))}dr^2 + r^2(d\Theta^2 + \sin^2\Theta d\phi^2),$$

the even modes $\xi^{(+)}_{2lm}(r, t)$ and odd modes $\xi^{(-)}_{2lm}(r, t)$ being built from a Fourier-type decomposition, i.e. [10]

$$\xi^{(+)}_{2lm}(r, t) = \int_{-\infty}^{\infty} dk \, a^{(+)}_{2klm}(r) \sin kt \frac{\sin kT}{\sin kT},$$

and

$$\xi^{(-)}_{2lm}(r, t) = \int_{-\infty}^{\infty} dk \, a^{(-)}_{2klm}(r) \cos kt \frac{\sin kT}{\sin kT},$$

where the radial functions $\xi^{(\pm)}_{2kl}$ obey the following second-order differential equation:

$$e^{-a} \frac{d}{dr} \left( e^{-a} \frac{d\xi^{(\pm)}_{2kl}}{dr} \right) + \left( k^2 - V^{\pm}_{l}(r) \right) \xi^{(\pm)}_{2kl} = 0,$$

where, on defining $\lambda \equiv \frac{(l+2)(l-1)}{2}$, the potential terms are given by [10]

$$V^{+}_{l}(r) = e^{-a(r,t)} \frac{2(\lambda^2(\lambda + 1)r^3 + 3\lambda^2mr^2 + 9)m^2r + 9m^3}{r^3(\lambda r + 3m)^2},$$

and

$$V^{-}_{l}(r) = e^{-a(r,t)} \left( \frac{l(l + 1)}{r^2} - \frac{6m}{r^3} \right),$$

respectively. In the expansion of the gravitational action to quadratic order, it is of crucial importance to evaluate the integral

$$I(k, k', l, R) \equiv \int_{0}^{R} e^{a(r,t)} \xi^{(+)}_{2kl}(r) \xi^{(+)}_{2k'l}(r) \, dr,$$

since [10] (see the appendix)
For this purpose, we bear in mind the limiting behaviours \[10\]

\[ \xi_{2kl}^{(\pm)} \sim \text{const} \times (kr)^{(l+1)} + O((kr)^{(l+3)}), \quad \text{as } r \to 0, \]

\[ \xi_{2kl}^{(\pm)}(r) \sim z_{2kl}^{(\pm)} e^{\pm kr}, \quad \text{as } r \to \infty, \]

where equation (16) results from imposing regularity at the origin, \( r_s \) is the Regge–Wheeler tortoise coordinate \[10, 16\]

\[ r_s(r) \equiv r + 2M \log(r - 2M), \]

\( z_{2kl}^{(\pm)} \) are complex constants. Indeed, it should be stressed that non-commutativity can smear plane waves into Gaussian wave packets. Thus, in a fully self-consistent analysis, the Fourier modes in equations (9) and (10), and their asymptotic form in equations (16) and (17), should be modified accordingly. However, this task goes beyond the aims of the present paper, and we hope to be able to perform these calculations in a future publication.

With this understanding, we can now exploit equation (11) to write the equations (hereafter, we write for simplicity of notation \( \xi_{kl} \) rather than \( \xi_{2kl}^{(\pm)} \), and similarly for \( V_l \) rather than \( V^{(\pm)}_l \))

\[ e^{a \xi_{kl}} \left[ e^{-a} \frac{d}{dr} \left( e^{-a} \frac{d}{dr} \xi_{kl} \right) + (k^2 - V_l) \xi_{kl} \right] = 0, \]

\[ e^{a \xi_{kl}} \left[ e^{-a} \frac{d}{dr} \left( e^{-a} \frac{d}{dr} \xi_{kl}' \right) + (k'^2 - V_l) \xi_{kl}' \right] = 0. \]

According to a standard procedure, if we subtract equation (20) from equation (19), and integrate the resulting equation from \( r = 0 \) to \( r = R \), we obtain

\[ (k^2 - k'^2) \int_0^R e^{a \xi_{kl}} \xi_{kl}' \, dr = \int_0^R \left[ \xi_{kl} \frac{d}{dr} \left( e^{-a} \frac{d}{dr} \xi_{kl} \right) - \xi_{kl}' \frac{d}{dr} \left( e^{-a} \frac{d}{dr} \xi_{kl} \right) \right] \, dr. \]

The desired integral (14) is therefore obtained from equation (21), whose right-hand side is then completely determined from the limiting behaviours in equations (16) and (17), i.e.

\[ \int_0^R e^{a \xi_{kl}} \xi_{kl}' \, dr = \left\{ \frac{1}{(k^2 - k'^2)} \left[ \xi_{kl} \frac{d}{dr} \left( e^{-a} \frac{d}{dr} \xi_{kl} \right) - \xi_{kl}' e^{-a} \frac{d}{dr} \xi_{kl} \right] \right\}_{r = 0}^{r = R}, \]

where, on going from equation (21) to equation (22), we have exploited the vanishing coefficient that weights the integral

\[ \int_0^R e^{-a} \left( \frac{d}{dr} \xi_{kl} \right) \left( \frac{d}{dr} \xi_{kl}' \right) \, dr, \]

resulting from two contributions of equal magnitude and the opposite sign. By virtue of equation (16), \( r = 0 \) gives vanishing contribution to the right-hand side of equation (22), while the contribution of first derivatives of radial functions involves also

\[ \frac{dr_s}{dr} \bigg|_{r = R} = \frac{R}{(R - 2M)}. \]

3. Generalized orthogonality relation

Note now that our metric (5) is a particular case of the spherically symmetric metric (8), since our \( a \) and \( b \) functions are independent of time. More precisely, unlike the full Vaidya spacetime, where in the region containing outgoing radiation the mass function varies extremely slowly
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with respect both to \( t \) and to \( r \) \cite{9}, we consider a ‘hybrid’ scheme where the mass function depends on \( r \) only for any fixed value of the non-commutativity parameter \( \theta \). We can therefore write

\[
e^{-a} = 1 - \frac{4M}{r \sqrt{\pi}} \left( \frac{3}{2} \right)^{\frac{1}{4}} \left( \frac{r}{4\theta} \right)
\]

in our non-commutative spherically symmetric model, where the function in curly brackets in equation (22) reads

\[
\left( 1 - \frac{4M}{R \sqrt{\pi}} \left( \frac{3}{2} \right)^{\frac{1}{4}} \left( \frac{R}{4\theta} \right) \right) \frac{R}{(R - 2M)} \left( \frac{1}{(k^2 - k'^2)} \right) \times \bar{e}^{i(k - k')_{cl} e^{i(k + k')_{cr}}(R)}
\]

\[
+ (k - k') z_{kl}^* z_{k'l'}^* e^{-i(k - k')_{cr}(R)} - (k + k') z_{kl} z_{k'l'}^* e^{i(k - k')_{cr}(R)}
\]

\[
+ (k + k') z_{kl}^* z_{k'l'} e^{i(k - k)_{cr}(R)} \right].
\]

At this stage, we exploit one of the familiar limits that can be used to express the Dirac \( \delta \), i.e. \cite{10}

\[
limit_{r_s \to \infty} e^{i(k \pm k')_{cr}} = i \pi \delta(k \pm k'),
\]

to find

\[
\int_0^R e^{i \xi_{kl} x^l_{k'} \, dr = 2\pi |z_{kl}|^2 \left( \delta(k + k') + \delta(k - k') \right),
\]

having defined

\[
F(R, \theta) = \frac{R}{(R - 2M)} \left[ 1 - \frac{4M}{R \sqrt{\pi}} \left( \frac{3}{2} \right)^{\frac{1}{4}} \left( \frac{R}{4\theta} \right) \right].
\]

4. Effect of \( \theta \) and expansion of the action functional

Since \( \theta \) has dimension length squared as we said after equation (3), we can define the non-commutativity-induced length scale

\[
L \equiv 2\sqrt{\theta}.
\]

Moreover, we know that our results only hold in the adiabatic approximation, i.e. when both \( m' \) and \( \dot{m} \) are very small. The latter condition is obviously satisfied because our mass function in equation (7) is independent of time. The former amounts to requiring that (hereafter we set \( w \equiv R/L \), while \( R_s \equiv 2M \))

\[
m'(R) = \frac{2}{\sqrt{\pi}} \frac{R_s}{L} e^{-w^2} w^2 \ll 1.
\]

Condition (29) is satisfied provided that either

(i) \( w \to \infty \) or \( w \to 0 \), i.e. \( R \gg L \) or \( R \ll L \);
(ii) or at \( R = L \) such that

\[
m'(R = L) = m'(w = 1) = \frac{2}{\sqrt{\pi}} \frac{R_s}{L} e^{-1} \ll 1,
\]

and hence for \( \frac{R_s}{L} \ll \frac{\sqrt{\pi}}{2} \).
Furthermore, at finite values of the non-commutativity parameter $\theta$, our $w \equiv R/L$ is always much larger than 1 in equation (27) if $R$ is very large, and hence we can exploit the asymptotic expansion of the lower incomplete $\gamma$-function in this limit [17, 18], i.e.

$$\gamma \left( \frac{3}{2}, w^2 \right) = \Gamma \left( \frac{3}{2} \right) - \Gamma \left( \frac{3}{2}, w^2 \right) \sim \frac{1}{2} \sqrt{\pi} \left[ 1 - e^{-w^2} \sum_{p=0}^{\infty} \frac{w^{1-2p}}{\Gamma \left( \frac{3}{2} - p \right)} \right]. \quad (31)$$

By virtue of equations (27) and (31), we find

$$F(R, \theta) \equiv F(R, L) \sim 1 + \frac{R_s}{(R - R_s)} e^{-w^2} \sum_{p=0}^{\infty} \frac{w^{1-2p}}{\Gamma \left( \frac{3}{2} - p \right)} \quad (32)$$

Equation (32) describes the asymptotic expansion of the correction factor $F$ when $R \gg L$.

In the opposite regime, i.e. for $\theta$ so large that $(R/L) \ll 1$ despite that $R$ tends to $\infty$, one has [18]

$$F(R, L) \sim \frac{R}{R_s} \left[ 1 - \frac{4}{3\sqrt{\pi}} \frac{R_s}{R} w^3 \left( 1 - \frac{7}{5} w^2 \right) \right]. \quad (33)$$

Last, but not least, if $R$ and $L$ are comparable, the lower-incomplete $\gamma$-function in equation (27) cannot be expanded, and we find, bearing in mind that $R_s/L \ll 1$ from equation (30), the limiting form

$$F(R, L) \sim 1 + R_s \frac{L}{R} \left( 1 - \frac{2}{\sqrt{\pi}} \gamma \left( \frac{3}{2}, 1 \right) \right) + O \left( \frac{(R_s/L)^2}{R} \right). \quad (34)$$

We therefore conclude that a $\theta$-dependent correction to the general relativity analysis in [10] does indeed arise from non-commutative geometry. In particular, the expansion of the action to quadratic order in perturbative modes takes the form (cf [10])

$$S_{(2)}^{\text{class}} = \frac{F(R, L)}{16} \sum_{l=2}^{\infty} \sum_{m=-l}^{l} \sum_{p=0}^{l} \frac{\Gamma(l-2)!}{(l+2)!} \int_0^{\infty} dk \, |z_{2l+m}|^2 |a_{2l+m} + P a_{-2l+m}|^2 \cot kT,$$

where the function $F(R, L)$ (see equation (27)) takes the limiting forms (32) and (33), respectively, depending on whether $w \gg 1$ or $w \ll 1$, while $P = \pm 1$ for even (respectively odd) metric perturbations. Our ‘correction’ $F(R, L)$ to the general relativity analysis is non-vanishing provided that one works at very large but finite values of $R$. In the limit as $R \to \infty$, one has instead

$$\lim_{R \to \infty} F(R, L) = 1, \quad (36)$$

which means that, at infinite distance from the Lorentzian singularity of Schwarzschild geometry, one cannot detect the effect of a non-commutativity parameter.

5. Absorption and emission spectra of a non-commutative Schwarzschild black hole

If general relativity is taken as the fundamental description of gravitational phenomena, particle emission by a black hole follows a Planck-type spectrum [19] while wave absorption by a black hole shows an oscillatory behaviour as a function of frequency [20]. For a Schwarzschild black hole, the two spectra are related by

$$H(\omega) = \frac{\sigma(\omega)}{\omega^{\frac{3}{2}}} \frac{T_H}{\pi} \frac{1}{1 - \frac{T_H}{\omega}}, \quad (37)$$

where $\omega$ is the frequency of the wave and $T_H$ is the temperature for a body emitting thermal
radiation, i.e. for the black hole. With this notation, \( H(\omega) \) and \( \sigma(\omega) \) are the emission and absorption rates, respectively. In [20] the author has found the following expression for the total absorption cross-section \( \sigma(\omega) \) in the Hawking formula (37):

\[
\sigma(\omega) = 27\pi M^2 - 2\sqrt{2}M \frac{\sin(2\sqrt{2\pi} \omega M)}{\omega},
\]

where \( \sigma(\omega) \) oscillates about the value \( 27\pi M^2 \) with decreasing amplitude \( \frac{2\sqrt{2}M}{\omega} \) and an approximately constant oscillation period.

We can ‘import’ the above result in our case by virtue of the adiabatic approximation exploited in [2–10], and we hence replace \( TH \) in equation (37) by the non-commutative black hole temperature \( T_{\text{HNC}} \) defined as

\[
T_{\text{HNC}} \equiv -\left( \frac{1}{4\pi} \frac{\partial g_{00}}{\partial r} \right)_{r=r_+} = \frac{1}{4\pi r_+} \left[ 1 - \frac{r_+^3}{4\theta^{3/2}} \frac{e^{-r_+^2/4\theta}}{\gamma(3/2; r_+^2/4\theta)} \right],
\]

where \( r_+ \) solves the horizon equation

\[
r_+ = 2m(r_+, \theta) \equiv \frac{4M}{\sqrt{\pi} \gamma \left( \frac{3}{2}, \frac{r_+^2}{4\theta} \right)}.
\]

Equation (38) yields therefore

\[
\sigma(\omega) = \frac{27\pi}{4} r_+^2 - \sqrt{2}r_+ \frac{\sin(\sqrt{2\pi} \omega r_+)}{\omega},
\]

where \( \sigma(\omega) \) oscillates about \( \frac{27\pi}{4} r_+^2 \) with decreasing amplitude \( \frac{\sqrt{2}r_+}{\omega} \) and a smaller oscillation period depending on \( \theta \) (note that the \( \theta \)-dependence in the horizon radius dominates at small \( \theta \)). Moreover, the emission spectrum of the black hole takes the form

\[
H(\omega) = \frac{\sigma(\omega)}{e^\omega / T_{\text{HNC}} - 1}.
\]

In the ‘large radius’ regime \( r_+^2/4\theta \gg 1 \), i.e. \( R \gg L \), equation (40) can be solved by iteration [14]. To first order in \( M/\sqrt{\theta} \), one finds

\[
r_+ = 2M \left( 1 - \frac{M}{\sqrt{\pi} \theta} e^{-M^2/4\theta} \right).
\]

The resulting plot of emission spectrum of the black hole (42) in the ‘large radius’ regime, i.e. if \( R \gg L \) \( (M \gg L) \) is shown in figure 1. In figure 1 the non-commutativity parameter in equation (28) is unable to modify the general relativity shape, although the maximum value attained is smaller than in general relativity.
6. Concluding remarks

Our paper has investigated the effect of non-commutative geometry on the recent theoretical analysis of quantum amplitudes in black hole evaporation, following the work in [1–10] (for other developments, see, e.g., the recent work in [21–24]). For this purpose, we have considered an approximate scheme where the background spacetime is static and spherically symmetric, with a mass function depending on the radial coordinate only for any fixed value of the non-commutativity parameter $\theta$.

Within this framework, we find that the general relativity analysis of spin-2 amplitudes is modified by a multiplicative factor $F$ defined in equation (27). Its limiting forms for $R \gg L$ or $R \ll L$ or $R \simeq L$ are given by equations (32), (33) and (34), respectively. Within this framework, unitarity is preserved, and the end state of black hole evaporation is a combination of outgoing radiation states (see section 1).

When the adiabatic approximation here assumed holds, we have also considered approximate formulae for the absorption and emission spectra of a ‘non-commutative Schwarzschild’ black hole. The resulting plot shows that for the total emission spectrum the general relativity shape is essentially recovered at finite values of $\theta$ such that $R \gg L$. Within such a scheme, Hawking emission is only important in a certain frequency range [20], but nothing can be said about the end-state of black hole evaporation.

An outstanding open problem is whether one can derive a time-dependent spherically symmetric background metric which incorporates the effects of non-commutative geometry. This would make it possible to improve the present comparison with the results in [2–10], where the Vaidya spacetime was taken as the background geometry. A closer inspection of the effect of a variable surface gravity is also in order, jointly with an assessment of the whole subject of black hole thermodynamics, when ‘corrected’ by non-commutative geometry.

Acknowledgments

The work of G Esposito and G Miele has been partially supported by PRIN SINTESI and PRIN FISICA ASTROPARTICELLARE, respectively.

Appendix A. From Regge–Wheeler to asymptotically flat gauge

Although we refer the reader to [10] for the large amount of detailed calculations, we should stress a few important points about the derivation of equation (15). The metric perturbations of linearized theory are subject to infinitesimal diffeomorphisms, which are the ‘gauge transformations’ of general relativity [25] (since the infinite-dimensional invariance group of the Einstein theory is indeed the diffeomorphism group).

The Regge–Wheeler (hereafter RW) gauge is not a supplementary condition, but rather an infinitesimal diffeomorphism according to which the odd-parity metric perturbations are modified as follows:

\begin{align*}
  h_{00}^{(-)\text{RW}} &= h_{00}^{(-)} - \partial_0 \Lambda_{lm}, \quad (A.1) \\
  h_{1lm}^{(-)\text{RW}} &= h_{1lm}^{(-)} - \partial_1 \Lambda_{lm} + \frac{2\Lambda_{lm}}{r}, \quad (A.2) \\
  h_{2lm}^{(-)\text{RW}} &= 0 = h_{2lm}^{(-)} + 2\Lambda_{lm}, \quad (A.3)
\end{align*}

where $h_{00}^{(-)}$ occurs in the expansion of the shift vector, while $h_{1lm}^{(-)}$ and $h_{2lm}^{(-)}$ occur in the expansion of the odd-parity 3-metric perturbations [10].
Figure 2. Plot of the function \( y = m'' + 2m'e^a(m' + rm)r^{-2} \), with \( m \) defined as in equation (7) and \( \theta = 1 \) (hence much smaller than \( R^2 \)).

In our spherically symmetric model of non-commutative gravity, since we rely upon the work in [14], where the left-hand side of the Einstein equations retains the same functional form as in general relativity, it remains possible to consider infinitesimal diffeomorphisms formally analogous to equations (A.1)–(A.3). For a more general framework, which we do not strictly need here, one should instead build a deformation of the algebra of diffeomorphisms, e.g. along the lines of the work in [26, 27].

In the RW gauge, one of the coupled partial differential equations relating \( h_{1lm} \) and \( h_{0lm} \) is [10]

\[
\partial_t^2 h_{1lm}^{\text{RW}} = \partial_t \partial_r h_{1lm}^{\text{RW}} - \frac{2}{r} \partial_t h_{0lm}^{\text{RW}} + \left[ -\frac{2\lambda e^b}{r^2} - \frac{2e^b}{r} \left( m'' + \frac{2m'e^a}{r^2}(m' + rm) \right) \right] h_{1lm}^{\text{RW}}. \tag{A.4}
\]

To recover the expected fall-off behaviour of metric perturbations, one later performs the asymptotically flat (hereafter AF) gauge transformation, according to which [10]

\[
h_{0lm}^{(-)AF} = 0 = h_{0lm}^{(-)RW} - \partial_t \Lambda_{lm}, \tag{A.5}
\]
\[
h_{1lm}^{(-)AF} = h_{1lm}^{(-)RW} - \partial_t \Lambda_{lm} + \frac{2\Lambda_{lm}}{r}, \tag{A.6}
\]
\[
h_{2lm}^{(-)AF} = h_{2lm}^{(-)RW} + 2\Lambda_{lm} = 2\Lambda_{lm}. \tag{A.7}
\]

By virtue of equations (A.6) and (A.4), and neglecting all derivatives of the mass function (this remains legitimate for our mass function in equation (7), as is clear from figure 2), one finds

\[
\partial_t^2 h_{1lm}^{(-)AF} = -\frac{2\lambda e^b}{r^2} h_{1lm}^{(-)RW}, \tag{A.8}
\]

since exact cancellations occur of the coefficients of \( \partial_r \partial_t \Lambda_{lm} \) and \( \partial_t^2 \Lambda_{lm} \). Moreover, since \( h_{2lm}^{(-)RW} \) vanishes in the RW gauge (see equation (A.3)), one finds, for the Zerilli function [10, 28],

\[
Q_{1lm}^{(-)RW} \equiv \frac{e^{-a}}{r} \left( h_{1lm}^{(-)RW} + \frac{r^2}{2} \partial_t h_{2lm}^{(-)RW} \right) = \frac{e^{-a}}{r} h_{1lm}^{(-)RW}. \tag{A.9}
\]

With our metric (5), one has \( e^b = e^{-a} \), and hence

\[
\partial_t^2 h_{1lm}^{(-)AF} = -\frac{2\lambda}{r} Q_{1lm}^{(-)RW}, \tag{A.10}
\]

By virtue of equations (A.8) and (A.10), and neglecting all derivatives of the mass function (this remains legitimate for our mass function in equation (7), as is clear from figure 2), one finds

\[
\partial_t^2 h_{1lm}^{(-)AF} = -\frac{2\lambda e^b}{r^2} h_{1lm}^{(-)RW}, \tag{A.8}
\]

since exact cancellations occur of the coefficients of \( \partial_r \partial_t \Lambda_{lm} \) and \( \partial_t^2 \Lambda_{lm} \). Moreover, since \( h_{2lm}^{(-)RW} \) vanishes in the RW gauge (see equation (A.3)), one finds, for the Zerilli function [10, 28],

\[
Q_{1lm}^{(-)RW} \equiv \frac{e^{-a}}{r} \left( h_{1lm}^{(-)RW} + \frac{r^2}{2} \partial_t h_{2lm}^{(-)RW} \right) = \frac{e^{-a}}{r} h_{1lm}^{(-)RW}. \tag{A.9}
\]

With our metric (5), one has \( e^b = e^{-a} \), and hence

\[
\partial_t^2 h_{1lm}^{(-)AF} = -\frac{2\lambda}{r} Q_{1lm}^{(-)RW}, \tag{A.10}
\]
in complete analogy with the general relativity analysis in [10]. Equation (A.10) is in turn used to prove the desired fall-off property
\[
h_{\gamma\mu}^{(-)}(r, t) = \frac{2\lambda}{r} \int_{-\infty}^{\infty} dk \frac{d_{\gamma\mu}^{(-)}(k)}{k^2} Q_{kl}^{(-)RW} \sin kt \sin kT, \tag{A.11}
\]
after writing [10]
\[
Q_{\gamma\mu}^{(-)RW} = \int_{-\infty}^{\infty} dk d_{\gamma\mu}^{(-)}(k) Q_{kl}^{(-)RW} \frac{\sin kt}{\sin kT}. \tag{A.12}
\]

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