Title: Beyond isotropic & homogeneous loop quantum cosmology: theory and predictions

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Abstract: In this talk I will discuss several recent advances in loop quantum cosmology and its extension to inhomogeneous models. I will focus on spherically symmetric spacetimes and Gowdy cosmologies with local rotational symmetry in vacuum. I will discuss how to implement a quantum Hamiltonian evolution on these quantizations. Then, I will focus on how we can extract predictions from those quantum geometries, and finally analyze a concrete example: cosmological perturbations on Bianchi I spacetimes in LQC.

Beyond isotropic & homogeneous loop quantum cosmology: theory and predictions

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Motivation

- Singularity theorems in classical GR.
- What is the nature of the spacetime close to the high curvature regions?
- Can these theories resolve all singularities?
- What are the observables we can measure?
- Can simple models provide predictions so that we can falsify them or even the full theory?

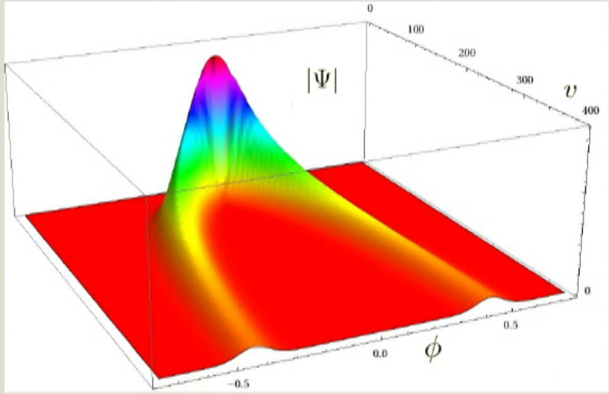
Motivation

- Symmetry reduced models allow us to realize concrete calculations (FRW, Bianchi, Schwarszchild, Kerr, ...).
- In quantum gravity, it is not obvious how to reduce the full quantum theory.
- The quantization of symmetry reduced models of GR can give us hints about the physics and mathematics of the full theory.
- For instance, we can study semiclassical sectors in agreement with GR and how quantum geometry can affect the predictions of the classical theory and its comparison with observations.

Homogeneous models in LQC

- Quantization of a FRW spacetime with a massless scalar field
  - a) Quantum dynamics (improved scheme) ✓
  - b) Singularity resolution (discrete quantum geometry) ✓
  - c) Semiclassical sectors and effective dynamics ✓

$$H^2 = \frac{8\pi G}{3}\rho\left(1 - \frac{\rho}{\rho_c}\right)$$



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- Extensions:

a) FRW with $k = 1$ ✓
b) FRW with $\Lambda = \pm 1$ ✓
c) Radiation dominated ✓
d) Bianchi I ✓
e) Bianchi II and IX
f) Kantowski-Sachs.

Homogeneous models in LQC

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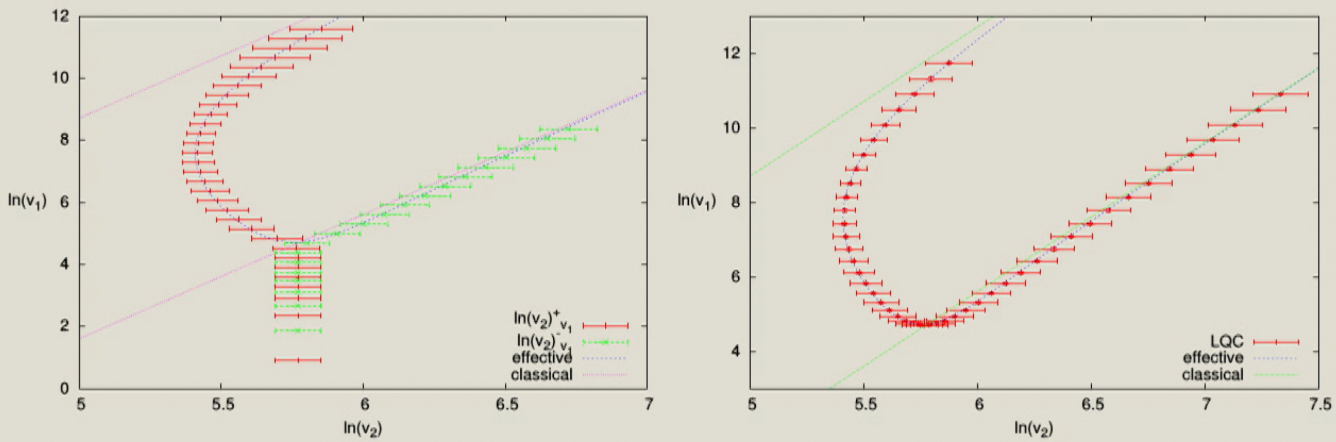


Martín-Benito, Mena Marugán, Pawlowski, Phys. Rev. D 80, 084038 (2009)

Inhomogeneous models

- Spherically symmetric vacuum spacetimes (Bojowald, Swiderski, 2004-2005) and polarized Gowdy models (Banerjee, Date, 2007).
  a) Symmetry reduction of the full theory in Ashtekar-Barbero variables.
  b) Robust kinematical quantum description.
  c) Well defined quantum Hamiltonian constraint (á la loop).
- Polarized Gowdy models – hybrid quantization (Martín-Benito, Mena-Marugán, Garay, 2008).
  a) Partial gauge fixing.
  b) Robust kinematical quantum description combining loop and Fock representation.
  c) Well defined quantum constraints (á la hybrid).
  d) Effective dynamics.

Inhomogeneous models

- Abelian constraint: Spherically symmetric spacetimes (Gambini, Pullin 2013; Gambini, O, Pullin 2013), (LRS) Gowdy cosmologies (Martín-de Blas, O, Pawlowski, 2015-2017) or 1+1 spacetimes like dilatonic black holes (Corichi, O, Rastgoo, 2016)
  a) classical abelianization of the Hamiltonian constraint,

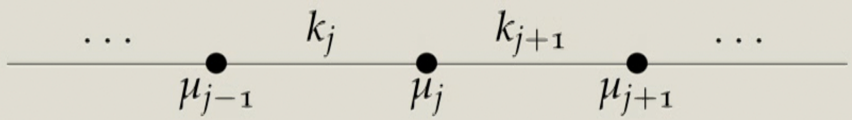
$$H_T = \int dx(NH + N^r H_r) = \int dx(\bar{N} H_{\rm new} + \bar{N}^r H_r),$$

$$H \to H_{\rm new} := \frac{(E^x)'}{E^\varphi} H - 2\frac{\sqrt{E^x}}{E^\varphi} K_\varphi H_r \Rightarrow \{H_{\rm new}(\bar{N}), H_{\rm new}(\bar{M})\} = 0.$$

Inhomogeneous models

- Abelian constraint: Spherically symmetric spacetimes (Gambini, Pullin 2013; Gambini, O, Pullin 2013), (LRS) Gowdy cosmologies (Martín-de Blas, O, Pawlowski, 2017) or 1+1 spacetimes like dilatonic black holes (Corichi, O, Rastgoo, 2016)
  - a) classical abelianization of the Hamiltonian constraint,
  - b) well-known kinematical Hilbert space,

$$|g, \vec{k}, \vec{\mu}\rangle =$$



$$\Psi_{g,\vec{k},\vec{\mu}}(A_x, A_\varphi) = \prod_{e_j \in g} \exp\left(i k_j \int_{e_j} dx\, A_x(x) \right) \prod_{v_j \in g} \exp\left(i \mu_j A_\varphi(v_j) \right).$$

$$\langle g', \vec{k}', \vec{\mu}' | g, \vec{k}, \vec{\mu}\rangle = \delta_{g',g}\delta_{\vec{k}',\vec{k}}\delta_{\vec{\mu}',\vec{\mu}}.$$

Inhomogeneous models

- Abelian constraint: Spherically symmetric spacetimes (Gambini, Pullin 2013; Gambini, O, Pullin 2013), (LRS) Gowdy cosmologies (Martín-de Blas, O, Pawlowski, 2017) or 1+1 spacetimes like dilatonic black holes (Corichi, O, Rastgoo, 2016)
  a) classical abelianization of the hamiltonian constraint
  b) well-known kinematical Hilbert space
  c) Physical Hilbert space

$$\|\Psi_{phys}\|^2 = \left(\langle\Psi_{kin}|\eta_{diff(H_r)}\eta_{diff(H_{\rm new})}\right)|\Psi_{kin}\rangle$$

  and Dirac observables

$$\hat{O}(z) = \ell_{\rm Pl}^2 \hat{k}_{{\rm Int}(nz)},$$

  and one physical global degree of freedom (either the mass $\hat{M}$ or a densitized shear scalar $\hat{h}$), are known.

Black hole spacetimes

- The effective spacetime metric can be computed by means of parametrized (relational) observables. For instance, in the spatially flat gauge $ds^2 = d\tilde{s}^2 + \langle \hat{E}^x(x) \rangle d^2\Omega$, where

$$d\tilde{s}^2 = -\langle \left(1 - \frac{\hat{r}_S}{\sqrt{\hat{E}^x(x)}} \right) \rangle dt^2 - \langle \eta \sqrt{\frac{\hat{r}_S}{[\hat{E}^x(x)]^{3/2}}} [\hat{E}^x(x)]' \rangle dt dx + \langle \left(\frac{[\hat{E}^x(x)]'}{4\hat{E}^x(x)} \right)^2 \rangle dx^2$$

- These geometries are discrete (piecewise constant x-functions).
- Close to where the singularity would be effective geometries are regular (singularity free).
- Interplay between the (fluctuating) discrete geometry and the (fluctuating) horizon (?).

Gowdy cosmologies with local rotational symmetry

- We also showed that it is possible to implement an improved dynamics scheme with $E^x \to V = \sqrt{\mathcal{E}} E^x$.
- On the physical Hilbert space, the quantum Hamiltonian evolution is defined as follows
  - Choice of phase space variable as time function T_j (on v_j).
  - A family of unitary (norm preserving) transformations $\hat{P}_{T_j}$ between $\mathcal{H}_{\text{phy}}$ and $\mathcal{H}_T$.
  - Relevant operators (observables) $\{\hat{O}(T)\}$ with suitable domains in $\mathcal{H}_T$.
- Finally, the evolution is defined via a family of operators $|\psi_{T'}\rangle := \hat{U}_{T',T}|\psi_T\rangle = \hat{P}_{T'}(\hat{P}_T)^{-1}|\psi_T\rangle$.
- As an example, we choose $V(\theta)$ as time function (non monotonic). Evolution split on several charts (on each vertex) (Martín-de Blas, O, Pawlowski - 2017).

Gowdy cosmologies with local rotational symmetry

- The solutions to the constraint and inner physical product can be written as

$$\tilde{\Psi}(\vec{k},h,\vec{v}) = \tilde{\Psi}(\vec{k},h)e_{\vec{k},h}(\vec{v}), \quad \langle\Phi|\Psi\rangle = \sum_{\vec{k}\in(\mathbb{Z}^\star)^n}\int_0^{h_\star(\vec{k})} dh\tilde{\Phi}^\star(\vec{k},h)\tilde{\Psi}(\vec{k},h).$$

- We split the solutions as via $\Psi^\pm(v) = \mathcal{F}^{-1}\theta(\pm b)[\mathcal{F}\Psi](b)$. This defines the unitary maps $\hat{U}^\pm_{\vec{v}',\vec{v}} = \hat{P}^\pm_{\vec{v}'}(\hat{P}^\pm_{\vec{v}})^{-1}$ with

$$\tilde{\Psi}^\pm_{\vec{v}}(\vec{k},h) = \hat{P}^\pm_{\vec{v}}\tilde{\Psi}(\vec{k},h) = \frac{e^\pm_{\vec{k},h}(\vec{v})}{|e^\pm_{\vec{k},h}(\vec{v})|}\tilde{\Psi}(\vec{k},h).$$

- The evolution on each chart, namely volume expanding and contracting, (in the Schrödinger picture)

$$|\Psi^\pm_{\vec{v}}\rangle = \sum_{\vec{k}\in(\mathbb{Z}^\star)^n}\int_0^{h_\star(\vec{k})} dh\tilde{\Psi}(\vec{k},h)\frac{e^\pm_{\vec{k},h}(\vec{v})}{|e^\pm_{\vec{k},h}(\vec{v})|}|\vec{k},h\rangle, \quad \langle\Psi^\pm_{\vec{v}}|\Psi^\pm_{\vec{v}}\rangle = 1, \quad \forall\vec{v}.$$

- Dirac observables $\hat{O}_{\vec{v}'} = \hat{U}^\pm_{\vec{v}',\vec{v}}\hat{O}_{\vec{v}}(\hat{U}^\pm_{\vec{v}',\vec{v}})^{-1}$

Gowdy cosmologies with local rotational symmetry

- Merits:
  - Rigorous quantum Hamiltonian (local) evolution.
  - It is not necessary to rely on the classical theory.
- Limitations:
  - The time-dependent states are not solutions to the constraint (but are related with them via a bijection).
  - The observables have ambiguous physical meaning close to the turning points (there it is more convenient to switch the time function).
  - Not obvious application to phase space variables with multiple turning points.
- Prospects: Extension of this Hamiltonian evolution in the context of black hole spacetimes.

Quantum field theories on quantum geometries

- Test fields (no backreaction) on these quantum geometries (perturbations) experience a dressed effective geometry.

$$\int d^4x \mathcal{L}_\phi^{\text{class}} = \frac{1}{2}\int dtd^3x \sqrt{-g}\left[-g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi - m_\phi^2\phi^2\right] \rightarrow$$
$$\int d^4x \hat{\mathcal{L}}_\phi^{\text{dress}} = \frac{1}{2}\int dtd^3x \left[-\langle\sqrt{-\hat{g}}\hat{g}^{\mu\nu}\rangle\partial_\mu\hat{\phi}\partial_\nu\hat{\phi} - \langle\sqrt{-\hat{g}}\rangle m_\phi^2\hat{\phi}^2\right]$$

- The propagation of fields on these (inhomogeneous) quantum geometries is not well understood.

- Extraction of predictions: role of discrete geometry and quantum anisotropies on the CMB; Hawking radiation, spectrum of quasi-normal modes (O - 2018), ...

QFT on quantum BHs

- For instance, the effective Hamiltonian of a test field (with parametrized observables in spatially flat gauge coordinates)

$$\hat{\mathcal{H}}_{\psi}^{\rm sph} = \langle\psi|\hat{\mathcal{H}}^{\rm sph}|\psi\rangle = \int dx \sum_{\ell m} \frac{1}{2}\Big(\mathcal{P}_1 \hat{\pi}^*_{\ell m}\hat{\pi}_{\ell m} + \mathcal{P}_2 \partial_x \hat{\phi}^*_{\ell m}\partial_x \hat{\phi}_{\ell m}$$

$$+ \left(\mathcal{P}_3 m_\phi^2 + \mathcal{P}_4 \ell(\ell+1)\right)\hat{\phi}^*_{\ell m}\hat{\phi}_{\ell m}\Big) + \hat{\pi}^*_{\ell m}\mathcal{P}_5 \partial_x \hat{\phi}_{\ell m},$$

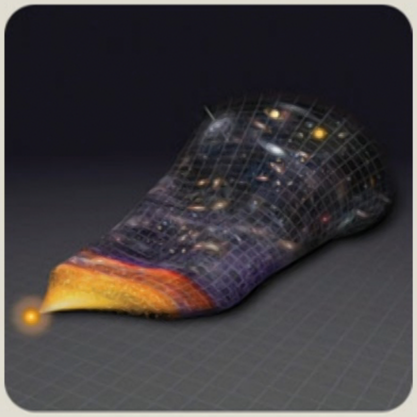
$$\mathcal{P}_1(x) = \langle\psi|\left[\widehat{\frac{2}{(E^x)'\sqrt{E^x}}}\right]|\psi\rangle, \quad \mathcal{P}_2(x) = \langle\psi|\left[\widehat{\frac{2\sqrt{E^x}E^x}{(E^x)'}}\right]|\psi\rangle, \quad \mathcal{P}_3(x) = \langle\psi|\left[\widehat{\frac{(E^x)'\sqrt{E^x}}{2}}\right]|\psi\rangle$$

$$\mathcal{P}_4(x) = \langle\psi|\left[\widehat{\frac{(E^x)'}{2\sqrt{E^x}}}\right]|\psi\rangle, \quad \mathcal{P}_5(x) = \langle\psi| - 2\eta\sqrt{2\hat{M}}\left[\widehat{\frac{(E^x)^{1/4}}{(E^x)'}}\right]|\psi\rangle.$$

$$\tilde{N}^2 = \mathcal{P}_4\sqrt{\mathcal{P}_1\mathcal{P}_2}, \quad \tilde{N}^x = \mathcal{P}_5, \quad \tilde{q}_{xx} = \frac{\mathcal{P}_4}{\sqrt{\mathcal{P}_1\mathcal{P}_2}}, \quad \tilde{q}_{\theta\theta} = \sqrt{\frac{\mathcal{P}_2}{\mathcal{P}_1}} = \frac{\tilde{q}_{\varphi\varphi}}{\sin^2\theta}, \quad \tilde{m}_\phi^2 = m_\phi^2\sqrt{\frac{\mathcal{P}_1}{\mathcal{P}_2}\frac{\mathcal{P}_3}{\mathcal{P}_4}}.$$

Example: Bianchi I & perturbations

- Homogeneity, isotropy and flatness are assumed in most of the present models of our Universe: inflation homogenizes and suppresses anisotropies and spatial curvature quickly.
- Together with quantum fluctuations of matter and geometry (quantum cosmological perturbations) we can provide predictions in very good agreement with observations.
- Upper bounds on anisotropies (shear) suggest that our Universe is isotropic at present.
- But this is not the end of the story ...



Example: Bianchi I & perturbations

- Cosmological perturbation theory on Bianchi I spacetimes has been studied in some detail (Pereira, Pitrou, Uzan, 2007-2008)
  a) rigorous formulation in terms of gauge-invariant perturbations (Lagrangian formulation and SVT decomposition),
  b) scalar and tensor perturbations are coupled dynamically if anisotropies are present,
  c) although anisotropies do not induce dipoles, they "break" scale invariance, isotropy and introduce scalar-tensor and tensor-tensor cross-correlations.
- Questions not fully understood:
  a) exact Fock quantization for the perturbations,
  b) extension to the Planck regime where quantum gravity is relevant (LQC),
  c) bounds on the anisotropies (shear) since perturbations can keep memory if they are large before (or close to) inflation.
  d) Characterization of the cross-correlations.

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Bianchi I spacetimes

- Let us consider a Bianchi I spacetime

$$ds^2 = -N^2 d\tau^2 + a_1^2 dx_1^2 + a_2^2 dx_2^2 + a_3^2 dx_3^2.$$

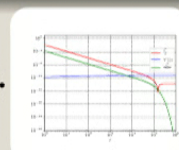
- We will assume a scalar field ϕ of mass m as matter content.

- The EOMs, in terms of $a = (a_1 a_2 a_3)^{1/3}$ are

$$3H^2 = \kappa\left[\frac{1}{2}\dot{\phi}^2 + V(\phi)\right] + \frac{1}{2}\tilde{\sigma}^2, \quad (\tilde{\sigma}^i_j)^{\cdot} = -3H\tilde{\sigma}^i_j,$$

$$\ddot{\phi} + 3H\dot{\phi} + V_\phi = 0, \qquad \tilde{\sigma}_{ij} = \frac{1}{2}\dot{\gamma}_{ij},$$

with $H = \frac{\dot{a}}{a}$, $\tilde{\sigma}^i_j = \gamma^{ik}\tilde{\sigma}_{kj}$, $\gamma_{ij} = \frac{h_{ij}}{a^2}$.



Bianchi I spacetimes

- Let us consider a Bianchi I spacetime
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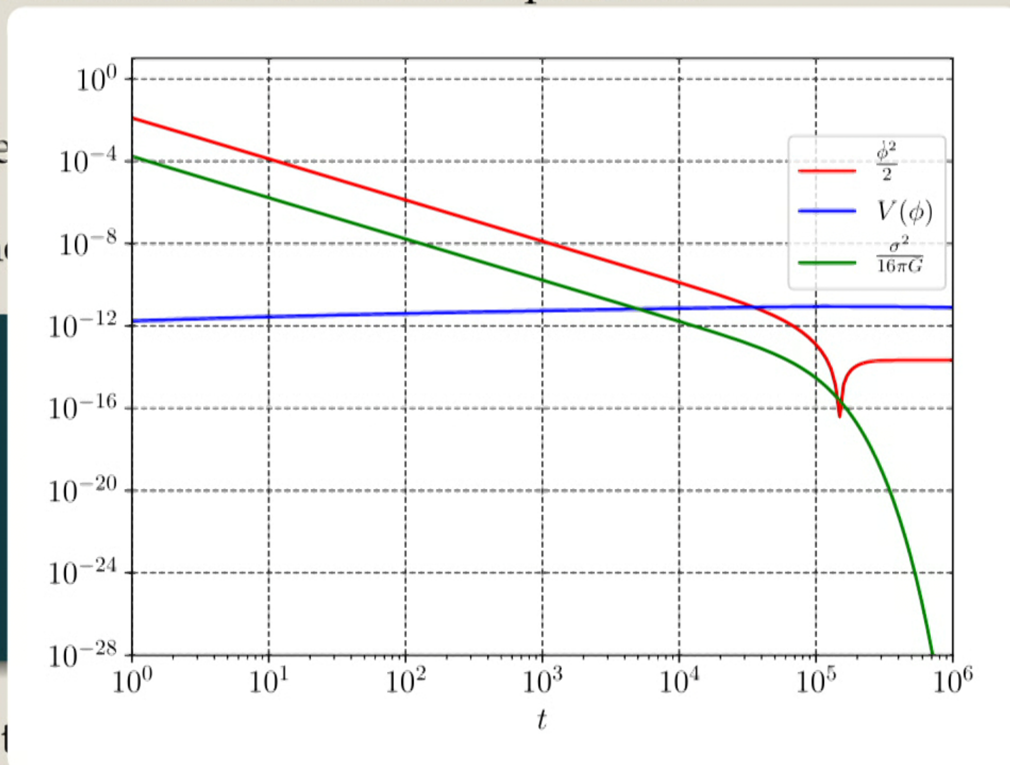


Perturbations: 2nd order Hamiltonian

- The 2nd order Hamiltonian takes the form

$$H^{(2)} = \frac{1}{2}\int d^3\mathbf{k}\left\{\pi_v\,\pi_v^* - \left[\frac{z_s''}{z_s} - k^2\right]v\,v^* - \sum_{\lambda=+,\times}\frac{1}{a^2}\left(2\sqrt{\kappa}\,\bar{\pi}_\phi\,\sigma_\lambda\mathcal{F}\right)'\left(v^*\,\mu_\lambda + v\,\mu_\lambda^*\right)\right.$$

$$\left. + \sum_{\lambda=+,\times}\pi_{\mu_\lambda}\,\pi_{\mu_\lambda}^* - \left[\frac{z_\lambda''}{z_\lambda} - k^2\right]\mu_\lambda\,\mu_\lambda^* + \left[2\,\sigma_+\,\sigma_\times - \left(2\,a^2\,\sigma_+\,\sigma_\times\mathcal{F}\right)'\right]\mu_{1-\lambda}\,\mu_\lambda^*\right\},$$

$$\mathcal{F} = \frac{2\,\mathcal{H} + \sigma_{||}}{\frac{32\pi G}{3}a^2\rho + \frac{2}{3}\left(2\,\sigma_{v_1}^2 + 2\,\sigma_{v_2}^2 + \sigma_+^2 + \sigma_\times^2\right)},$$

$$\frac{z_s''}{z_s} = \frac{a''}{a} - a^2\,V_{\phi\phi} + \frac{1}{a^2}\left(2\,\kappa\,\bar{\pi}_\phi^2\mathcal{F}\right)',$$

$$\frac{z_\lambda''}{z_\lambda} = \frac{a''}{a} + 2\,\sigma_{(1-\lambda)}^2 + \frac{1}{a^2}\left(a^2\,\sigma_{||}\right)' + \frac{1}{a^2}\left(2\,a^2\,\sigma_\lambda^2\mathcal{F}\right)'.$$

QFT for cosmological perturbations

- The Fourier modes $\gamma = (\vec{q}, \vec{\pi})$, with $\vec{q} = (v, \mu_+, \mu_\times)$ and $\vec{\pi}$ the conjugate momenta, can be written as $\vec{q} = \sum_{i=1}^{3} a_i \vec{q}^{(i)} + \text{c.c.}$, a linear combination of the (orthonormal) basis of complex solutions $\gamma^{(i)}$ and $(\gamma^{(i)})^*$ with respect to the norm

$$\langle \gamma, \tilde{\gamma} \rangle = i\hbar \sum_{j=1}^{3} \left(q_j^\star \tilde{\pi}^j - (\pi^j)^\star \tilde{q}_j \right).$$

Note: addends "j" preserved independently if fields uncoupled.

- Quantum fields given by

$$\hat{\vec{q}} = \sum_{i=1}^{3} \hat{a}_i \vec{q}^{(i)} + \hat{a}_j^\dagger \left(\vec{q}^{(i)} \right)^*, \quad \left[\hat{a}_i(\boldsymbol{k}), \hat{a}_j^\dagger(\boldsymbol{k}') \right] = \delta_{ij} \delta^{(3)}(\boldsymbol{k} + \boldsymbol{k}')$$

Bianchi I spacetimes in LQC

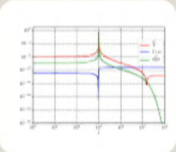
- In LQC the effective dynamics is determined by (Ashtekar, Wilson-Ewing, Mena-Marugán, Martín-Benito, ...)

$$\mathcal{H}_{\rm eff} = \frac{1}{8\,\pi G\,\gamma^2\,v}\left[-\left(\frac{\sin(\bar{\mu}_1\,c_1)}{\bar{\mu}_1}\,\frac{\sin(\bar{\mu}_2\,c_2)}{\bar{\mu}_2}p_1p_2 + \text{cyclic terms}\right) + \bar{\pi}_\phi^2 + v^2V(\phi)\right],$$

where $v^2 = p_1p_2p_3$, $\bar{\mu}_1 = \lambda\sqrt{\frac{|p_1|}{|p_2\,p_3|}}$ (and similarly for $\bar{\mu}_2$ and $\bar{\mu}_3$), and $\lambda^2 = 4\sqrt{3}\,\pi\,\gamma\,\ell_{\rm Pl}^2$.

- The energy density, mean Hubble parameter and shear are bounded above (Gupt, Singh, 2012-2013)

$$\rho_{\rm max} = 0.41\rho_{\rm Pl}, \quad H_{\rm max} = \frac{8.34}{\ell_{\rm Pl}}, \quad \sigma^2_{\rm max} = \frac{11.57}{\ell_{\rm Pl}^2}.$$



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Bianchi I spacetimes in LQC

- In I … r, Wilson-Ewing,
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$\mathcal{H}_{ef}$ … $+ v^2 V(\phi)\Big]$,

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λ^2 = …

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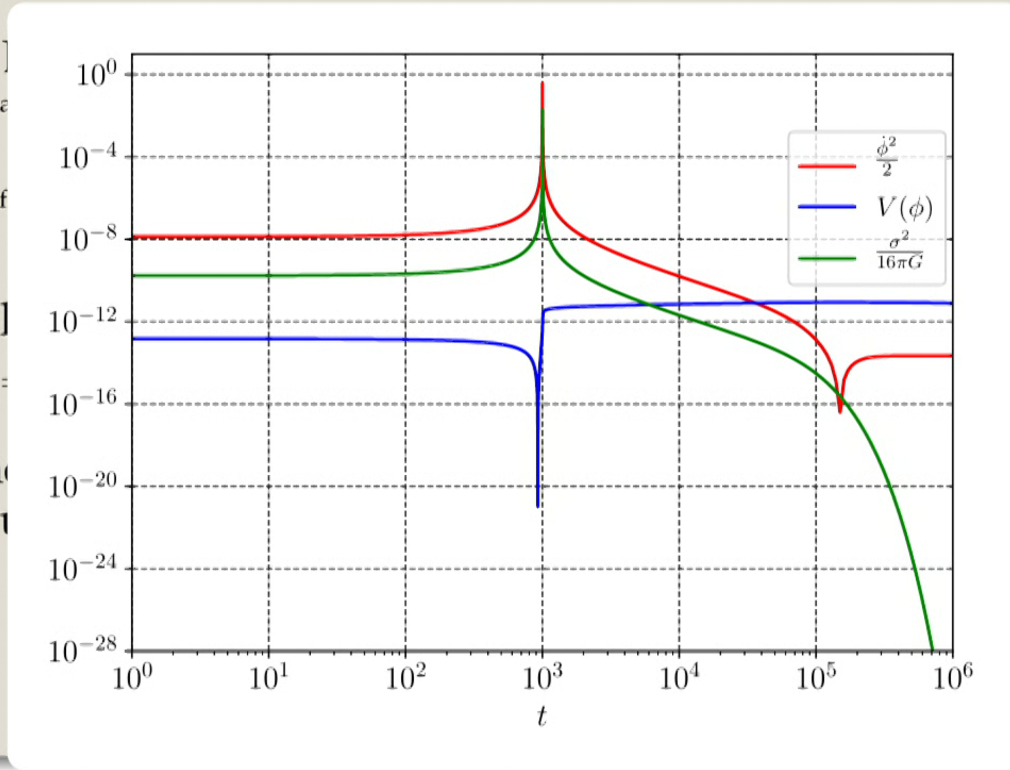
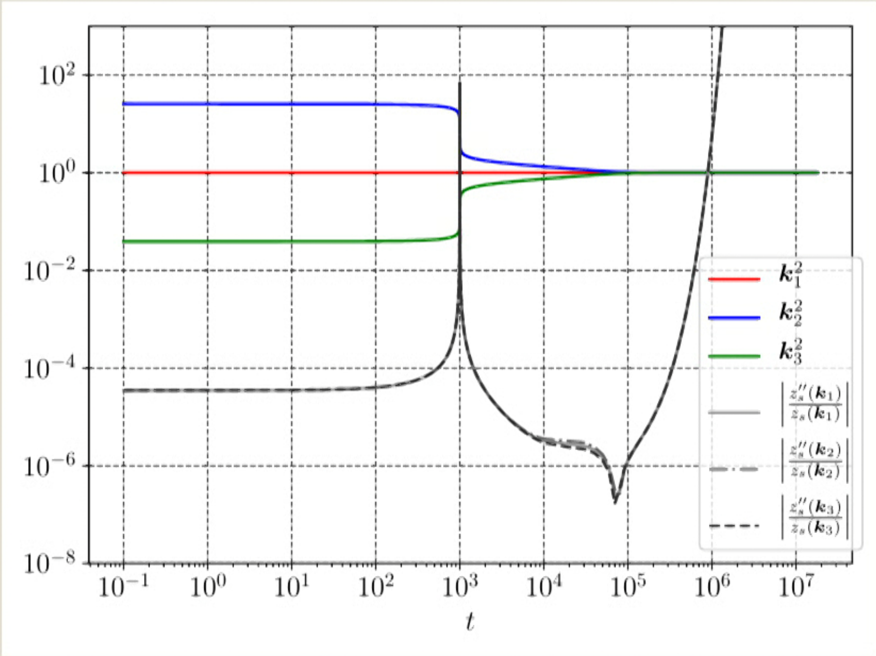
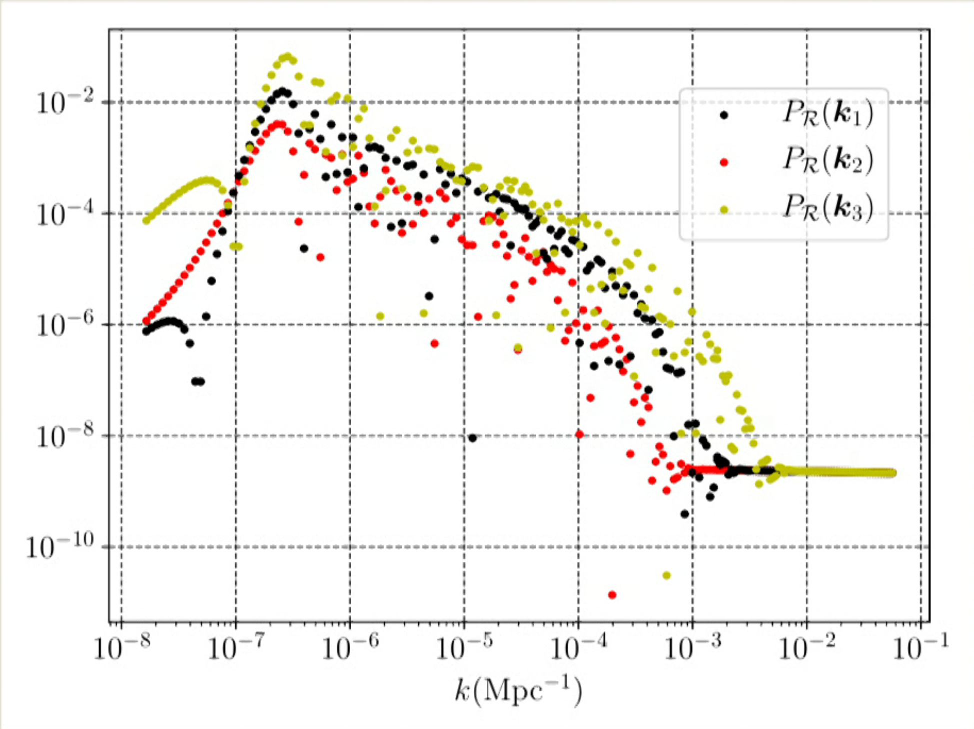


Adiabatic conditions



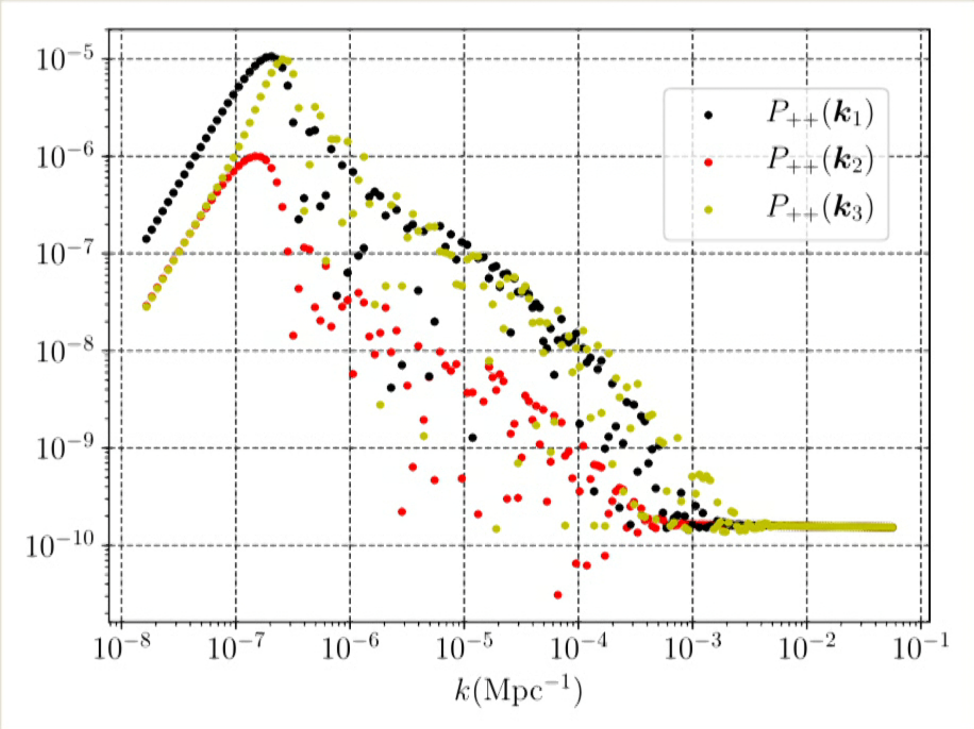
Before the bounce the spacetime is (approximately) isotropic and there is a well defined adiabatic regime where we assume the 0th order adiabatic vacuum state.

Scalar power spectra



The isotropy of the scalar and tensor power spectra at the end of inflation is broken for large wevalengths.

Tensor power spectra



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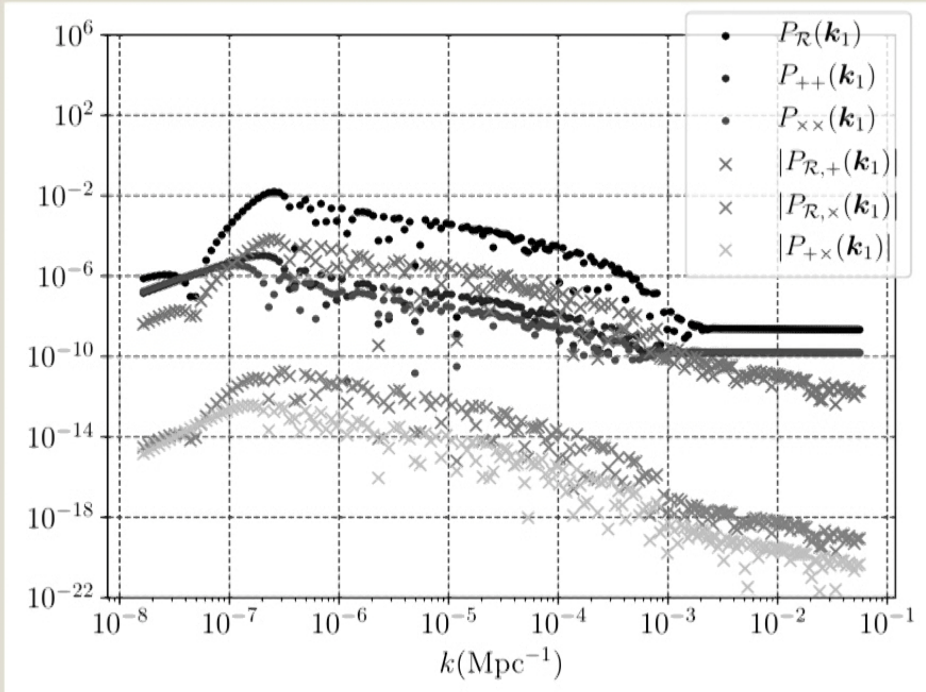
Correlation functions



Cross-correlations can be large for small wavenumbers. For large wavenumbers they behave as $P_{AB} = \frac{P_{AB}^{(0)}(\hat{\boldsymbol{k}})}{k} \cos(\phi_{AB}(\boldsymbol{k}))$, with $(A, B) = \{\mathcal{R}, +, \times\}$ and where $P_{AB}^{(0)}(\hat{\boldsymbol{k}})$ grow with $\sigma^2(t_B)$.

Upper bounds in the shear

- Given a fixed numer of e-folds, it is possible to find upper bounds for the shear. For $k = k'$, let us define the estimator

$$\Delta\mathcal{P}_{\mathcal{R}}(\mathbf{k},\mathbf{k}',\delta k) = 2\left|\frac{\bar{\mathcal{P}}_{\mathcal{R}}(\mathbf{k},\delta k) - \bar{\mathcal{P}}_{\mathcal{R}}(\mathbf{k}',\delta k)}{\bar{\mathcal{P}}_{\mathcal{R}}(\mathbf{k},\delta k) + \bar{\mathcal{P}}_{\mathcal{R}}(\mathbf{k}',\delta k)}\right|.$$

Then, for $\phi_B = 1.083$ (i.e. k_{LQC} at $\ell \simeq 20$), and for $k = k' = 2.1 \cdot 10^{-3}(Mpc)^{-1}$ (or equivalently $\ell \simeq 30$), we have

$\sigma^2(t_B)$	$\Delta\mathcal{P}_{\mathcal{R}}(\mathbf{k}_1,\mathbf{k}_2,\delta k)$	$\Delta\mathcal{P}_{\mathcal{R}}(\mathbf{k}_1,\mathbf{k}_3,\delta k)$	$\Delta\mathcal{P}_{\mathcal{R}}(\mathbf{k}_2,\mathbf{k}_3,\delta k)$
0.1	0.038	0.31	0.36
0.2	0.024	0.71	0.73
0.3	0.038	1.0	1.1
0.4	0.033	1.25	1.27
1.0	0.020	1.84	1.85

Summary

- Loop quantum gravity techniques applied to symmetry reduced models successfully deals with several fundamental questions: singularity resolution, semiclassical geometries, effective description, etc.

- Full quantization of inhomogeneous cosmological and black hole models is available.

- Their quantum dynamics, semiclassical sectors and effective geometries need to be further explored.

- Test quantum fields on these quantum geometries deserve additional attention (dressed effective metric): the fundamental discretization and fluctuations of these quantum geometries can modify our understanding of several well-known phenomena of QFTs on classical curved spacetimes. Further research in this direction seems very promising.