Title: Graviton propagator from EPRL spinfoam model

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Abstract: We derive geometric correlation functions in the new spinfoam model with coherent states techniques, making connection with quantum Regge calculus and perturbative quantum gravity. In particular we recover the expected scaling with distance for all components of the propagator. We expect the same technique to be well-suited for other spinfoam models.

Graviton propagator from new spinfoam models

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Università degli studi Roma Tre

Perimeter Institute 2009

Harmonic-radial gauge in classical EM and GR

CP, Magliaro, Rovelli

Wave-packet propagation in QG

CP, Magliaro, Rovelli - CP, Alesci, Bianchi, Magliaro

Asymptotics of fusion coefficients

CP, Alesci, Bianchi, Magliaro

Divergencies (?) of new spinfoam models

CP, Rovelli, Speziale

Graviton propagator in the new models

CP, Magliaro

Outline of the talk

- Graviton propagator - the physical picture

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- Spinfoams
- New models

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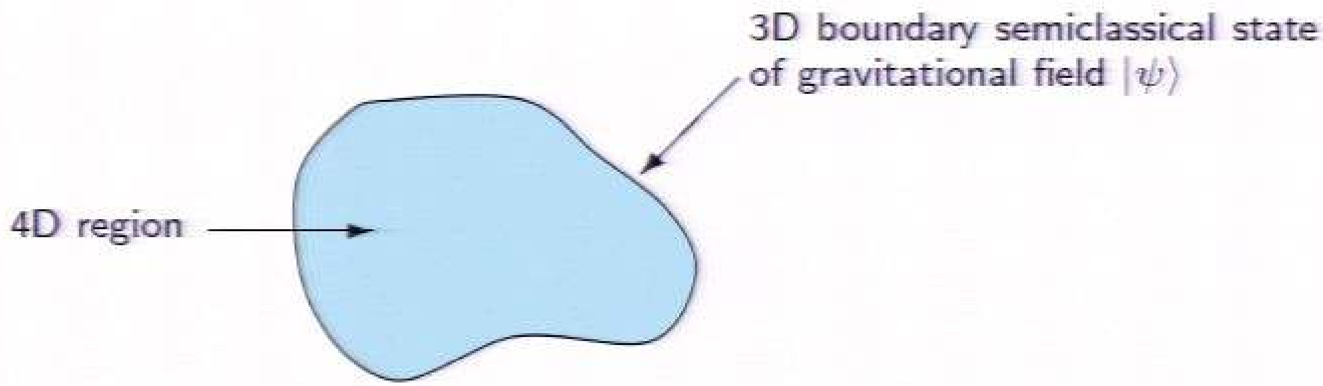
- Spinfoams
- New models
- The new propagator
- Conclusions, work in progress and outlook

Graviton propagator - the physical picture Colosi, Conrady, Doplicher Modesto, Oeckl, Rovelli, Testa



Spinfoams assign an amplitude to this boundary state: $\langle W|\psi\rangle$

$$\text{If } \psi = \psi_{\text{in}}^{*} \otimes \psi_{\text{out}} \rightarrow \langle W|\psi\rangle \sim \langle \psi_{\text{out}}| e^{iHt} |\psi_{\text{in}}\rangle$$

$\langle W|$ codes the dynamics of quantum general relativity.

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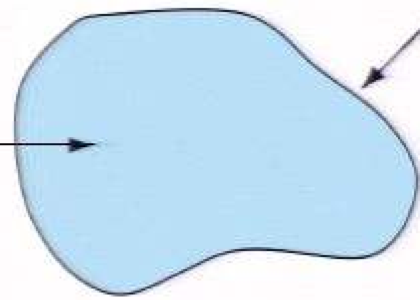

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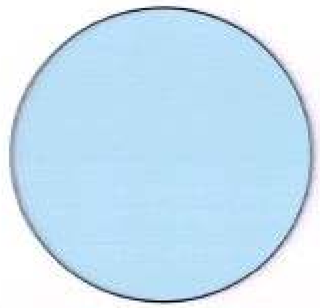
$\langle W|$ codes the dynamics of quantum general relativity. Define

$$\langle\cdot\rangle = \frac{\langle W|\cdot|\psi\rangle}{\langle W|\psi\rangle}$$

that we call physical expectation value, in contrast to the kinematical expectation value

$$\langle\cdot\rangle_{\text{kin}} = \frac{\langle\psi|\cdot|\psi\rangle_{\text{kin}}}{\langle\psi|\psi\rangle_{\text{kin}}}$$

Graviton propagator - the physical picture



boundary semiclassical state coding the geometry (intrinsic and extrinsic) of a Euclidean 3-sphere

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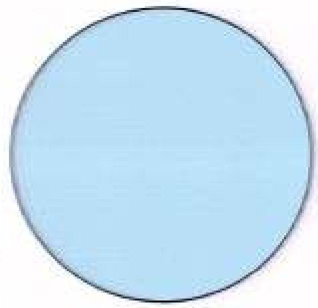

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Then construct an excited boundary state by acting with metric field operators at two points:

$$(g(x) - < g(x) >)(g(y) - < g(y) >)\,|\psi\rangle$$

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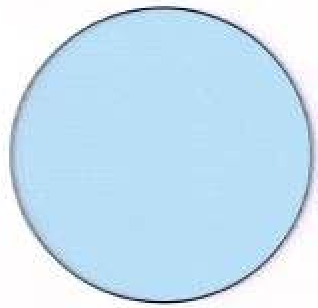

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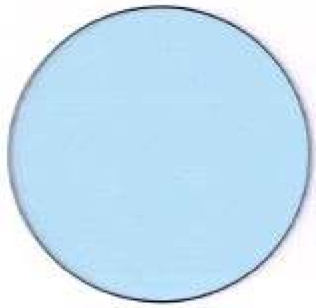
It is called 2-point function, or graviton propagator, in (nonperturbative) quantum gravity. *What is it*? It represents the probability of detecting two excitations at the spacetime points x and y over flat space-time (here we consider only the Euclidean signature). More correctly, by flat space-time we mean a semiclassical state peaked over this classical geometry.

Graviton propagator - the physical picture



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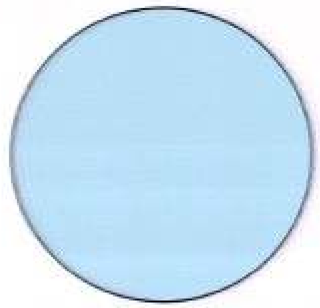

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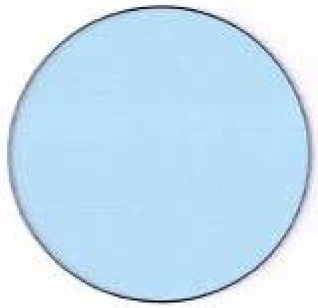

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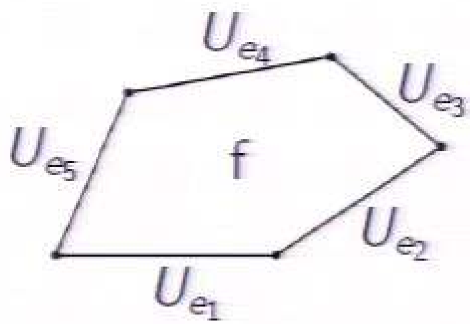
Perturbativity inside non-perturbativity!

Spinfoams - discretization of classical

Start with a triangulated manifold (faces/triangles, edges/tetrahedra, vertices/4-simplices) and dicretize the Holst-Plebanski action for GR:



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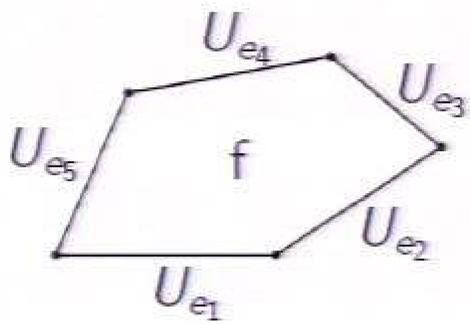

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- **Simplicity constraint** $*B_f \cdot B_f = 0$
- **Cross-simplicity constraint** $*B_f \cdot B_{f'} = 0 \quad f, f' \subset t$
- **Closure constraint** $\sum_{f \subset t} B_f = 0$

Replace cross-simplicity by:

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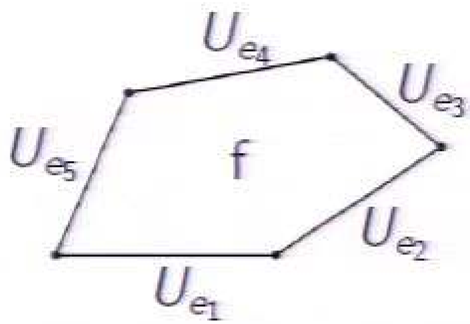

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Boundary variables: $B_f(t_1), B_f(t_2), U_f(t_1, t_2)$

The variable conjugate to U_f is $J_f = *B_f + \frac{1}{\gamma} B_f$

Spinfoams - quantization

Quantize the theory choosing an appropriate Hilbert space. Similarly to lattice Yang-Mills theory, define the kinematical Hilbert space as (G=Spin(4))

$$L^2(G^L) \qquad L = \# \text{ of boundary faces (links)}$$

and quantize

$$J_f(t_1) \longrightarrow \text{l.i. vector field}$$
$$J_f(t_2) \longrightarrow \text{r.i. vector field}$$

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- **Closure constraint** project onto the gauge-invariant subspace

$$f(g) = f(hgh^{-1})$$

  Then an orthonormal basis of this subspace is given by spin-networks where each link is labeled by irreps of Spin(4)=SU(2)⊗SU(2), i.e. couples of spins (j^+, j^-), and each node is labeled by intertwiners between irreps

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- **Simplicity constraint** restricts these spin labels to be of the form

$$j^+ = \left|\frac{1+\gamma}{1-\gamma}\right| j^-$$

Spinfoams - the model of Barrett and Crane

- **Strong cross-simplicity constraint** In the Barrett-Crane model the cross-simplicity constraint is imposed (too) strongly and restricts the intertwiner space to the 1-dimensional space spanned by the BC intertwiner. The vertex amplitude (4-simplex amplitude) depends *only* on the spins. Two problems arise:

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the boundary state space doesn't match with the one of LQG

and

despite some components of the graviton propagator work well (Bianchi, Modesto, Rovelli, Speziale), the other components have the wrong scaling (Alesci, Rovelli)

These problems could trace back to the fact that we have overconstrained the system; the intertwiner space is too small!

Spinfoams - new models: $EPRL_\gamma$ & FK_γ

Engle, Pereira, Rovelli, Livine
Freidel, Krasnov

Consider for simplicity the $EPRL_\gamma$ model ($=FK_\gamma$ for $\gamma < 1$)

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- **Weak cross-simplicity constraint** We can solve weakly this constraint. Alternatively we impose the single master constraint:

$$C^i \equiv B^{0i} = 0 \;\Leftrightarrow\; |\vec{C}| = 0 \quad \text{(master constraint)}$$

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Each (j^+, j^-) component in the Peter-Weyl decomposition of $L^2(\text{Spin}(4))$ can be further decomposed:

$$(j^+, j^-) = j^+ \otimes j^- = |j^+ - j^-| \oplus \ldots \oplus (j^+ + j^-)$$

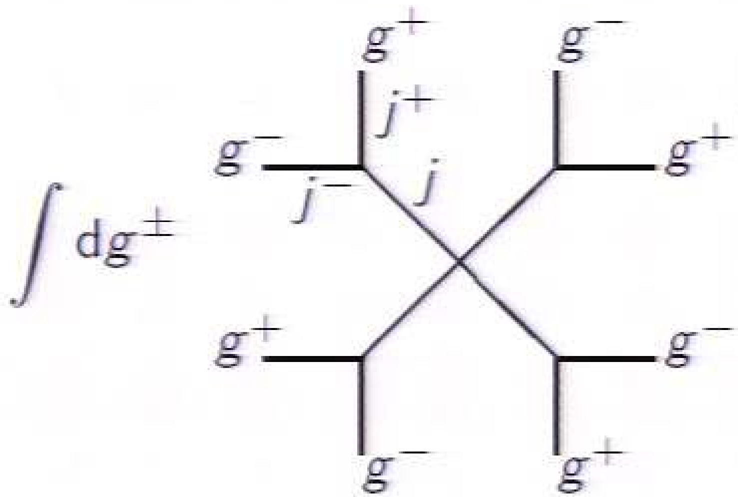
The cross-simplicity constraint selects the component j satisfying the following relation

$$j^\pm = \frac{|1 \pm \gamma|}{2} j$$

i.e. the highest or the lowest weight.

Spinfoams - new models (2)

$\Longrightarrow$ intertwiners are of the form ($1 \leftrightarrow 1$ with LQG intertwiner space):



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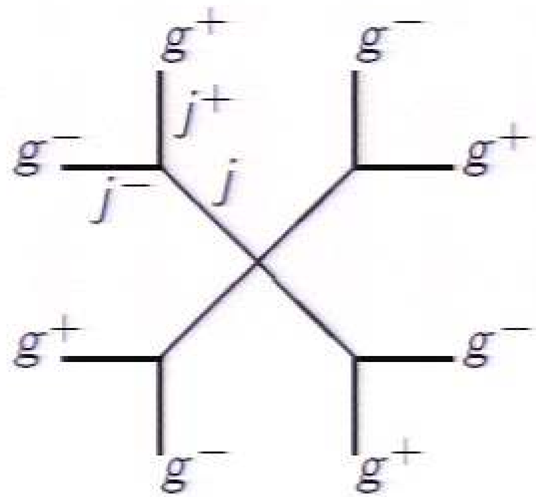
$\int dg^{\pm}$ 

The 4-simplex amplitude is the contraction of 5 of them. In the recoupling basis:

$$W_{\mathrm{EPRL}}(j,i) = \sum_{i_a^+, i_a^-} 15j(j^+, i^+) 15j(j^-, i^-) f(i_1, i_1^+, i_1^-) \ldots f(i_5, i_5^+, i_5^-)$$

Where

$$f(i, i^+, i^-) = i^{abcd} C_a^{a^+ a^-} \ldots C_d^{d^+ d^-} i^+_{a^+ b^+ c^+ d^+} i^-_{a^- b^- c^- d^-}$$

Coherent intertwiners

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The important property is

$$\mathbb{1}_j = \sum_m |j, m\rangle\langle j, m| = \int \mathrm{d}g |j, g\rangle\langle j, g| = \int \mathrm{d}n |j, n\rangle\langle j, n|$$

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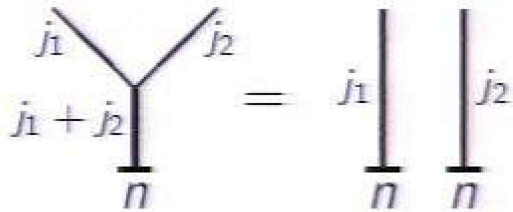
so that coherent states form an overcomplete basis of H_j. We define a coherent (4-valent) intertwiner by taking the tensor product of four coherent states and then projecting onto the invariant subspace:

$$|j_1 \dots j_4, n_1 \dots n_4\rangle = \int \mathrm{d}g \; g \triangleright |j_1, n_1\rangle \otimes \dots \otimes |j_4, n_4\rangle$$

Spinfoams - EPRL$_\gamma$ model (3)

Consider the EPRL$_{\gamma<1}$ or FK$_{\gamma<1}$ vertex in the coherent intertwiner basis. Using the following decomposition property of coherent states



we can untie the nodes (see 2) and rewrite the vertex amplitude as a product of contractions of coherent states:

$$W(j,n)=\int \mathrm{d}^5 g^{\pm} \prod_{a<b} \langle -n_{ab}|(g_a^{+})^{-1} g_b^{+}|n_{ba}\rangle^{2j_{ab}^{+}} \prod_{a<b} \langle -n_{ab}|(g_a^{-})^{-1} g_b^{-}|n_{ba}\rangle^{2j_{ab}^{-}}$$

We have used also the tensoring property of coherent states

$$|1/2,n\rangle \underbrace{\otimes \ldots \otimes}_{2j \text{ times}} |1/2,n\rangle = |j,n\rangle$$

to write in powers of the 1/2 representation.

Asymptotic analysis Barrett, Fairbairn, Dowdall, Gomes, Hellmann - Conrady, Freidel

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Consider this expression for large j's.

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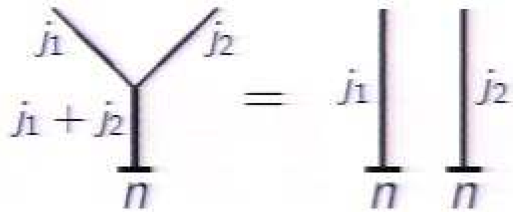

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$$\max \operatorname{Re} S \implies g^{\pm}_{a} n_{ab} = -g^{\pm}_{b} n_{ba}$$

$$\frac{\partial S}{\partial g} = 0 \implies \sum_{b \neq a} j_{ab} n_{ab} = 0$$

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It follows:

$$(g^{\pm}_{a})^{-1}g^{\pm}_{b}|n_{ba}\rangle = e^{i\phi^{\pm}_{ab}}|-n_{ab}\rangle \qquad \phi^{+}_{ab}-\phi^{-}_{ab}=\Phi_{ab} \text{ dihedral angle between } t_a \text{ and } t_b$$

Then

$$W(j,n)\sim\begin{cases} A(j)(e^{iS_{\text{Regge}}}+\text{c.c}) & n=n(j)\\ \text{suppressed} & \text{otherwise}\end{cases} \qquad S_{\text{Regge}}=\sum \gamma j_{ab}\Phi_{ab}(j)$$

The new propagator

$$G_{nm}^{abcd} = \langle \mathbf{E}_n^a \cdot \mathbf{E}_n^b \, \mathbf{E}_m^c \cdot \mathbf{E}_m^d \rangle - \langle \mathbf{E}_n^a \cdot \mathbf{E}_n^b \rangle \langle \mathbf{E}_m^c \cdot \mathbf{E}_m^d \rangle$$

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- New boundary state superposition of semiclassical coherent 4-simplices, peacked around $j_{ab} = j_0$, and $\delta j / j_0 \to 0$

$$|\Psi\rangle = \sum_j \psi(\mathbf{j}) \, |\mathbf{j}\rangle_0 \quad \psi(\mathbf{j}) \equiv e^{-\frac{1}{j_0} \delta \mathbf{j}^T \alpha \delta \mathbf{j} + i\gamma\phi_0 \sum_{a<b} \delta j_{ab}}$$

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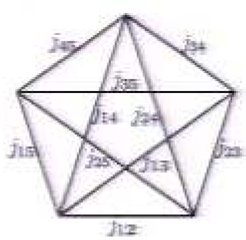

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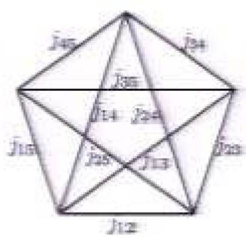

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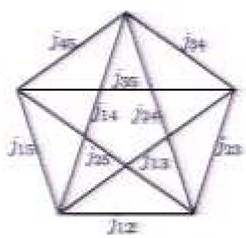

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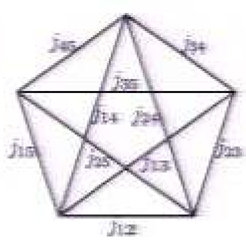

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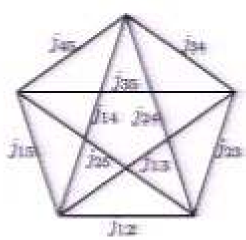

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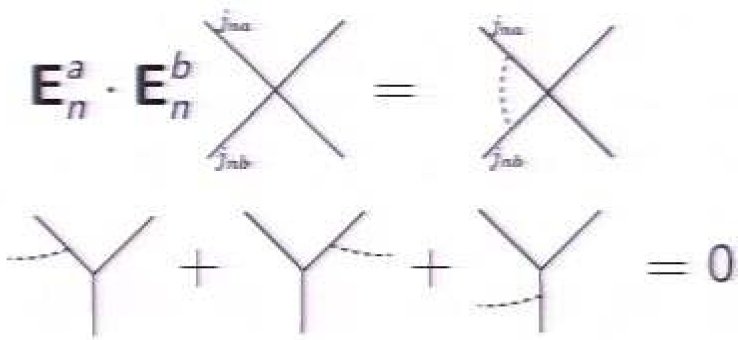

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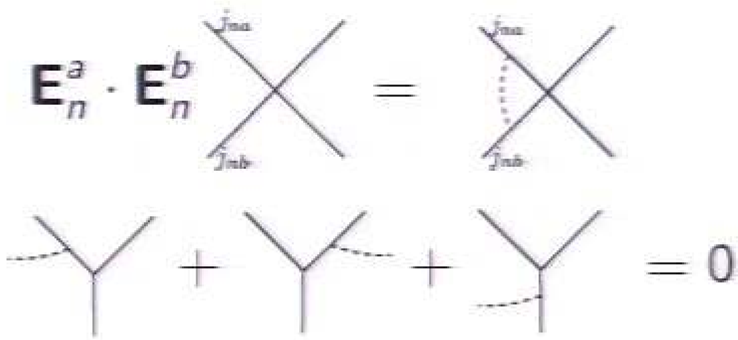

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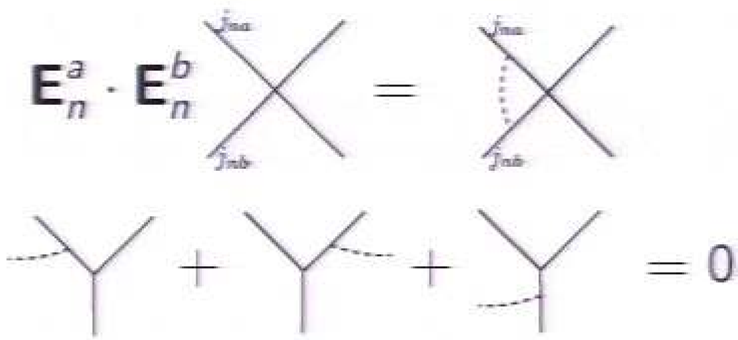

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The propagator is a simplicial non-perturbative version of the standard propagator of perturbative theory:

$$G_{nm}^{abcd} \longrightarrow G^{\mu\nu\rho\sigma}(x,y) = \langle 0 | h^{\mu\nu}(x) h^{\rho\sigma}(y) | 0 \rangle$$

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Consider the action on a coherent 4-simplex,

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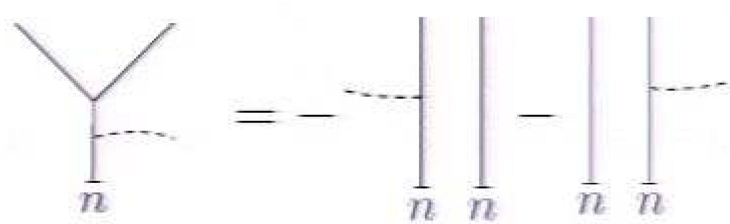

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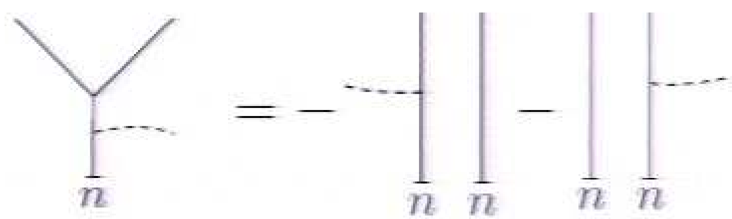

$$\int \mathrm{d}^5 g^{\pm} \langle j_{na}, -n_{na}|\sigma(g_n^+)^{-1} g_a^+ |j_{na}, n_{an}\rangle \langle j_{nb}, -n_{nb}|\sigma(g_n^+)^{-1} g_b^+ |j_{nb}, n_{bn}\rangle$$

$$\times \left(\text{non-grasped factors}\right) + \ldots$$

It is just the new vertex evaluated on a grasped-intertwiner configuration. This expression seems quite complicated *but*...

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…we have the identity:

$$\langle j, n_1|\sigma|j, n_2\rangle = j\frac{\langle 1/2, n_1|\sigma|1/2, n_2\rangle}{\langle 1/2, n_1|1/2, n_2\rangle}\langle j, n_1|j, n_2\rangle \equiv j\frac{\langle n_1|\sigma|n_2\rangle}{\langle n_1|n_2\rangle}\langle j, n_1|j, n_2\rangle$$

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$$\mathbf{A}_n^a|_{\text{saddle}} = j_{na}n_{na}$$

namely, on the saddle point $\mathbf{A}$ (that we racall we obtained recasting the action of the grasping operator as an insertion in the group integral) is the classical simplicial quantity measured by the grasping operator.

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All ingredients are ready. The propagator in terms of **A** is:

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$$G_{nm}^{aacd} \sim \partial_j (j_{na}^2) \partial_j (j_{mc} j_{md} n_{mc} \cdot n_{md}) \left(iS_{\mathrm{R}}'' - \frac{\alpha}{j^0} \right)^{-1}$$

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$$G_{nm}^{aacc} \sim \partial_j (j_{na}^2) \partial_j (j_{mc}^2) \left(iS_R'' - \frac{\alpha}{j^0} \right)^{-1} = 4 j_0^2 \left(iS_R'' - \frac{\alpha}{j^0} \right)^{-1}_{(na)(mc)}$$

- **Diagonal-nondiagonal components:** $a = b, c \neq d$ $\partial_g \mathbf{A}_n^a \cdot \mathbf{A}_n^a$ vanishes, while $\partial_g \mathbf{A}_m^c \cdot \mathbf{A}_m^d$ is different from zero. But since the jg block is zero, also for those components only the Regge sector contributes:

$$G_{nm}^{aacd} \sim \partial_j (j_{na}^2) \partial_j (j_{mc} j_{md} n_{mc} \cdot n_{md}) \left(iS_R'' - \frac{\alpha}{j^0} \right)^{-1}$$

- **Nondiagonal components:** $a \neq b, c \neq d$ Here we have one contribution from the Regge sector and one contribution from the non-Regge sector:

$$G_{nm}^{abcd} \sim \partial_j (j_{na} j_{nb} n_{na} \cdot n_{nb}) \partial_j (j_{mc} j_{md} n_{mc} \cdot n_{md}) \left(iS_R'' - \frac{\alpha}{j_0} \right)^{-1} + \text{non-Regge term}$$

The new propagator (7)

Final remarks

The new propagator has the right scaling properties. When all the normalizations are correctly implemented, it reproduces the Newtonian law

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for all components, contrary to the BC model, in which the sole diagonal-diagonal components had the correct scaling. In fact that was a major motivation for searching new (corrected) spinfoam models.

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We have computed the "non-Regge term" of nondiagonal-nondiagonal components. Though we didn't expect to find it, it seems to be there, and we don't have a clear physical interpretation of it. Nevertheless, it could be required in order to match the tensorial structure with the one of linearized theory. This is a topic I am working on.

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The new propagator (4)

...we have the identity:

$$\langle j, n_1|\sigma|j, n_2\rangle = j\frac{\langle 1/2, n_1|\sigma|1/2, n_2\rangle}{\langle 1/2, n_1|1/2, n_2\rangle}\langle j, n_1|j, n_2\rangle \equiv j\frac{\langle n_1|\sigma|n_2\rangle}{\langle n_1|n_2\rangle}\langle j, n_1|j, n_2\rangle$$

Using this, the (grasped) 4-simplex evaluation becomes simpler...

$$\langle W|\mathbf{E}_n^a \cdot \mathbf{E}_n^b|\mathbf{j}\rangle_0 = \int \mathrm{d}^5 g^{\pm}\, \mathbf{A}_n^a \cdot \mathbf{A}_n^b\; e^{S(g)}$$

where $\mathbf{A}_n^a = A_n^{a+} + A_n^{a-}$ $\qquad A_n^{a\pm} = j_{na}^{\pm}\frac{\langle -n_{na}|\sigma(g_n^{\pm})^{-1}g_a^{\pm}|n_{an}\rangle}{\langle -n_{na}|(g_n^{\pm})^{-1}g_a^{\pm}|n_{an}\rangle}$

Since we are interested in the large distance (large j_0) limit, we may regard $\mathbf{A}$ as an insertion and evaluate the integral for large spins with the saddle point method. Remarkably,

$$\mathbf{A}_n^a|_{\text{saddle}} = j_{na}n_{na}$$

namely, on the saddle point $\mathbf{A}$ (that we racall we obtained recasting the action of the grasping operator as an insertion in the group integral) is the classical simplicial quantity measured by the grasping operator.

The new propagator (5)

All ingredients are ready. The propagator in terms of **A** is:

$$\frac{\sum_j \psi(j) \int \mathrm{d}g^{\pm} \mathbf{A}\cdot\mathbf{A}\,\mathbf{A}\cdot\mathbf{A} e^S}{\sum_j \psi(j) \int \mathrm{d}g^{\pm} e^S} - \frac{\sum_j \psi(j) \int \mathrm{d}g^{\pm} \mathbf{A}\cdot\mathbf{A} e^S}{\sum_j \psi(j) \int \mathrm{d}g^{\pm} e^S} \frac{\sum_j \psi(j) \int \mathrm{d}g^{\pm} \mathbf{A}\cdot\mathbf{A} e^S}{\sum_j \psi(j) \int \mathrm{d}g^{\pm} e^S}$$

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- We should check what happens when going beyond the single 4-simplex level and/or switching to the Lorentzian signature.

We could...we could...we could...

We could...we could...we could...
That's why I love physics!

thank you for the attention

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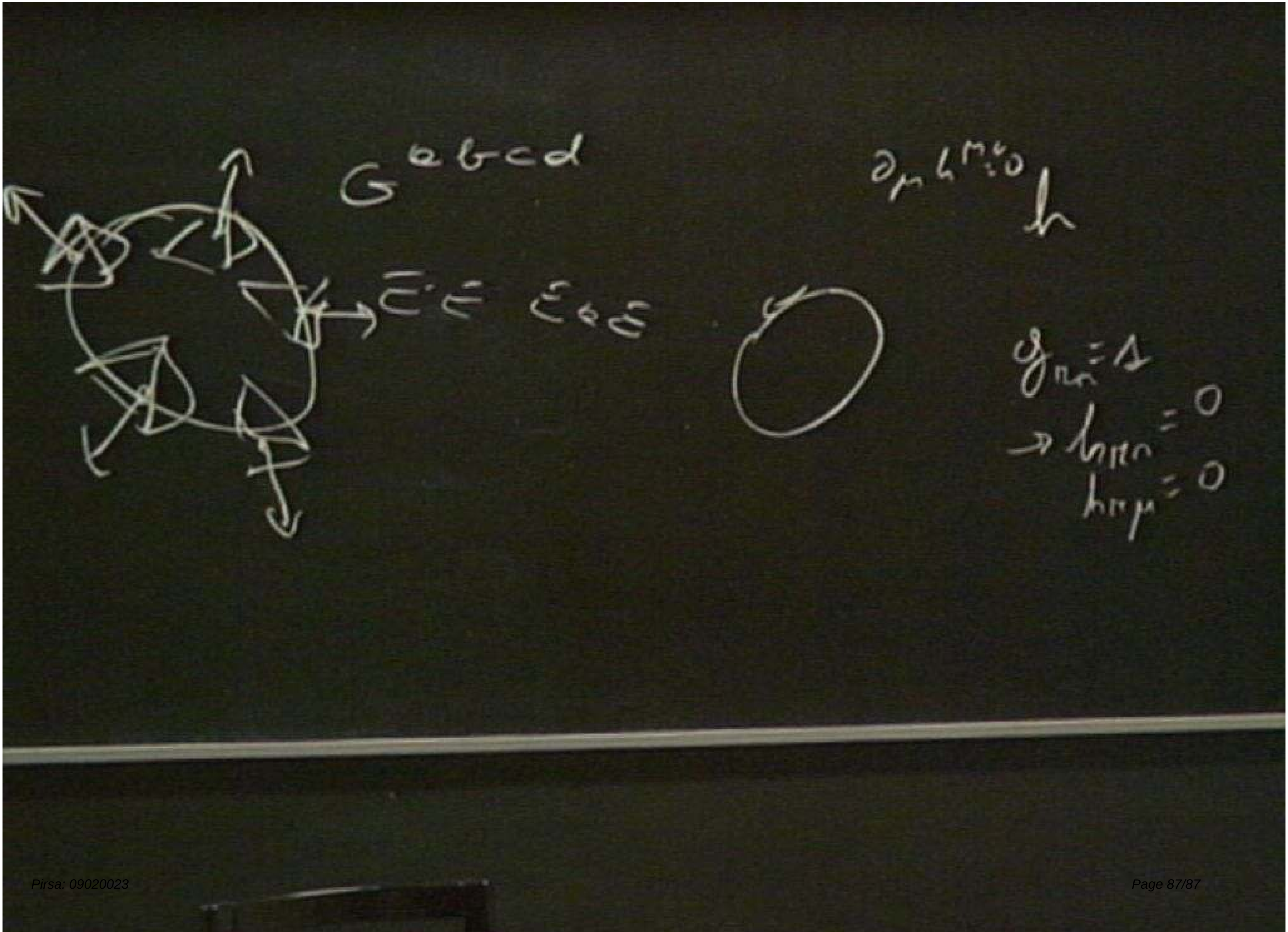
for long beautiful discussions together


Gabcd
∂μ h^μν=0
h
g_nn = 1
h_nn = 0
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