Construction of the $\Phi^4$-quantum field theory on noncommutative Moyal space

Harald Grosse$^1$ and Raimar Wulkenhaar$^2$

$^1$ Fakultät für Physik, Universität Wien
Boltzmanngasse 5, A-1090 Wien, Austria

$^2$ Mathematisches Institut der Westfälischen Wilhelms-Universität
Einsteinstraße 62, D-48149 Münster, Germany

Abstract

We review our recent construction of the $\phi^4$-model on four-dimensional Moyal space. A milestone is the exact solution of the quartic matrix model $Z[E,J] = \int d\Phi \exp(\text{trace}(J\Phi - E\Phi^2 - \frac{\lambda}{4!}\Phi^4))$ in terms of the solution of a non-linear equation for the 2-point function and the eigenvalues of $E$. The $\beta$-function vanishes identically. For the Moyal model, the theory of Carleman type singular integral equations reduces the construction to a fixed point problem. Its numerical solution reveals a second-order phase transition at $\lambda_c \approx -0.396$ and a phase transition of infinite order at $\lambda = 0$. The resulting Schwinger functions in position space are symmetric and invariant under the full Euclidean group. They are only sensitive to diagonal matrix correlation functions, and clustering is violated. The Schwinger 2-point function is reflection positive iff the diagonal matrix 2-point function is a Stieltjes function. Numerically this seems to be the case for coupling constants $\lambda \in [\lambda_c, 0]$.

1 Introduction

Perturbatively renormalised quantum field theory is an enormous phenomenological success, a success which lacks a mathematical understanding. The perturbation series is at best an asymptotic expansion which cannot converge at physical coupling constants. Some physical effects such as confinement are out of reach for perturbation theory. In two and partly three dimensions, methods of constructive physics [GJ87, Riv91], often combined with the Euclidean approach [Sch59, OS73, OS75], were used to rigorously establish quantum field theory models.

In four dimensions there was little success so far. It is generally believed that

---

$^*$based on the lectures given at RIMS, Kyoto University, September 2013

$^1$harald.grosse@univie.ac.at

$^2$raimar@math.uni-muenster.de
due to asymptotic freedom, non-Abelian gauge theory (i.e. Yang-Mills theory) has
the chance of a rigorous construction. But this is a hard problem [JW00]. What
makes it so difficult is the fact that any simpler model such as quantum electro-
dynamics or the $\lambda \phi^4$-model cannot be constructed in four dimensions (Landau
ghost problem [LAK54a, LAK54b, LAK54c] or triviality [Aiz81, Frö82]).

One of the main difficulties is the non-linearity of the models under consid-
eration. Fixed point methods provide a standard approach to non-linear problems,
but they are rarely used in quantum field theory. In this contribution we review
a sequence of papers [GW12b, GW13b, GW14] in which we successfully used
symmetry and fixed point methods to exactly solve a toy model for a quantum
field theory in four dimensions.

1. Following [GW12b], we show in sec. 2 that a Ward identity for the $U(\infty)$
group action leads to an exact solution of the quartic matrix model $Z =
\int D[\Phi] \exp(\text{trace}(J\Phi - E\Phi^2 - \frac{1}{4}\Phi^4))$ in terms of the solution of a non-linear
equation. As by-product we find that any renormalisable quartic matrix
model has vanishing $\beta$-function. All these steps are completely elementary.

2. Self-dual $\phi^4_4$-theory on Moyal space [GW05b, GW05c] is of that type. For
extreme noncommutativity $\theta \to \infty$, and after careful discussion of ther-
modynamic and continuum limit, the non-linear equation is reduced to
a fixed-point problem [GW12b] which has a unique non-perturbative and
non-trivial solution for $\lambda < 0$ [GW14]. Sec. 3 reviews this work. The key
step is the observation that a certain difference function satisfies a linear
singular integral equation of Carleman type [Car22, Tri57]. We also present
some numerical results, contained in work in progress [GW14], which show
evidence for phase transitions.

3. Following [GW13b], we identify in sec. 4 a limit to Schwinger functions
for a scalar field on $\mathbb{R}^4$. Surprisingly for a highly noncommutative model,
these Schwinger functions show full Euclidean symmetry. Otherwise they
have unusual properties such as absent momentum transfer in interaction
processes. This seems to suggest triviality, but the numerical investiga-
tion [GW14] of the 2-point function shows scattering remnants from a non-
commutative geometrical substructure. Most surprisingly, the Schwinger
2-point function seems to be reflection positive in one of its phases.

2 Exact solution of the quartic matrix model

For us a ‘matrix’ is a compact (Hilbert-Schmidt) operator on Hilbert space $H =
L^2(I, \mu)$. Such operators $\Phi \in L^2(H)$ can be represented by integral kernel oper-
ators $(\Phi v)_a = \int_I d\mu_b \Phi_{ab} v_b$. Then all natural matrix operations such as product,
adjoint and trace have counterparts $(\Phi \Phi')_{ab} = \int_I d\mu_c \Phi_{ac} \Phi'_{cb}$, $(\Phi^*)_{ab} = \overline{\Phi_{ba}}$ and
$\text{tr}(\Phi \Phi') = \int_I d\mu_a (\Phi \Phi')_{aa}$ in $L^2(H)$.
To define a Euclidean quantum field theory for a matrix $\Phi \in \mathcal{L}^2(H)$ we give ourselves an action functional

$$S[\Phi] = V \text{tr}(E\Phi^2 + P[\Phi]) .$$

(1)

Here, $P[\Phi]$ is a polynomial in $\Phi$ with scalar coefficients, and this alone would be a familiar action in the theory of matrix models [DGZ95]. To be closer to field theory on a (compact) manifold $M$ we add the analogue of the kinetic term $\int_M dx (-\Delta \phi) \phi$, that is, we require the external matrix $E$ to be an unbounded selfadjoint positive operator on $H$ with compact resolvent. The volume $V$ will play a crucial rôle. The construction involves several regularisation and limiting procedures. One such regularisation consists in a finite size $N$ for the matrices, and $V$ will be a certain function of $N$ which together with $N$ is sent to $\infty$.

Adding a source term to the action, we define the partition function as

$$Z[J] = \int \mathcal{D}[\Phi] \exp(-S[\Phi] + V \text{tr}(\Phi J)) ,$$

(2)

where $\mathcal{D}[\Phi]$ is the extension of the Lebesgue measure from finite-rank operators to $\mathcal{L}^2(H)$ and $J$ a test function matrix. For absent $P[\Phi] \mapsto 0$ in (1), $\mathcal{D}[\Phi] \exp(-V \text{tr}(E\Phi^2))/Z[0]$ would be the Gaussian measure of covariance determined by $E$. What we want, and what we achieve, is to construct (the moments of) the measure $\mathcal{D}[\Phi] \exp(-V \text{tr}(E\Phi^2 + \frac{\lambda}{4} \Phi^4))/Z[0]$ in the limit $V \to \infty$. Such a limit cannot be expected for $Z$. Instead, we pass to the generating functional $\log Z[J]$ of connected correlation functions,

$$\langle \varphi_{a_1b_1} \cdots \varphi_{a_NB_N} \rangle_c = \frac{\partial^N \log Z[J]}{\partial J_{b_1a_1} \cdots \partial J_{b_Na_N}} \bigg|_{J=0} .$$

(3)

2.1 Ward identity and topological expansion

Unitary operators $U$ belonging to an appropriate unitisation of the compact operators on $H$ give rise to a transformation $\Phi \mapsto \tilde{\Phi} = U\Phi U^*$. Since the space of selfadjoint compact operators is invariant under the adjoint action, we have

$$\int \mathcal{D}[\Phi] \exp(-S[\Phi] + V \text{tr}(\Phi J)) = \int \mathcal{D}[\tilde{\Phi}] \exp(-S[\tilde{\Phi}] + V \text{tr}(\tilde{\Phi} J)) .$$

Unitary invariance $\mathcal{D}[\tilde{\Phi}] = \mathcal{D}[\Phi]$ of the Lebesgue measure implies

$$0 = \int \mathcal{D}[\Phi] \left\{ \exp(-S[\Phi] + V \text{tr}(\Phi J)) - \exp(-S[\Phi] + V \text{tr}(\Phi J)) \right\} .$$

Note that the integrand $\cdots$ itself does not vanish because $\text{tr}(E\Phi^2)$ and $\text{tr}(\Phi J)$ are not unitarily invariant; we only have $\text{tr}(P[\Phi]) = \text{tr}(P[\tilde{\Phi}])$ due to $UU^* = U^* U = \text{id}$.
together with the trace property. Linearisation of $U$ about the identity operator leads to the Ward identity

$$0 = \int \mathcal{D}[\Phi] \left\{ E\Phi\Phi - \Phi\Phi E - J\Phi + \Phi J \right\} \exp(-S[\Phi] + V \text{tr}(\Phi J)) .$$

(4)

We can always place ourselves in an orthonormal basis of $H$ where $E$ is diagonal (but $J$ is not). Since $E$ is of compact resolvent, $E$ has eigenvalues $E_a > 0$ of finite multiplicity $\mu_a$. We thus label the matrices by an enumeration of the (necessarily discrete) eigenvalues of $E$ and an enumeration of the basis vectors of the finite-dimensional eigenspaces. Writing $\Phi$ in \{\ldots\} of (4) as functional derivative $\Phi_{ab} = \frac{\partial}{\partial \nu_{J_{ab}}}$, we have proved (first obtained in \cite{DGMR07}):

**Proposition 1** The partition function $Z[J]$ of the matrix model defined by the external matrix $E$ satisfies the $|I| \times |I|$ Ward identities

$$0 = \sum_{n \in I} \left( \frac{E_a - E_p}{V} \frac{\partial^2 Z}{\partial J_{an} \partial J_{np}} + J_{pn} \frac{\partial Z}{\partial J_{an}} - J_{na} \frac{\partial Z}{\partial J_{np}} \right) .$$

(5)

Without loss of generality we can assume that the map $I \ni m \mapsto E_m \in \mathbb{R}_+$ is injective. Namely, correlation functions will only depend on the set of eigenvalues ($E_m$) of $E$. Partitioning the index set $I$ into equivalence classes $[m]$ which have the same $E_m$, the index sum over a function that only depends on $E_m$ becomes $\sum_{m \in I} f(m) = \sum_{[m] \in [I]} \mu_{[m]} f([m])$. Therefore, at the expense of adding a measure $\mu_{[m]} = \dim \ker(E - E_m \text{id})$, we can assume that $m \mapsto E_m$ is injective.

In a perturbative expansion, Feynman graphs in matrix models are ribbon graphs. Viewed as simplicial complexes, they encode the topology $(B, g)$ of a genus-$g$ Riemann surface with $B$ boundary components (or punctures, marked points, holes, broken/external faces). Some simple examples for $P[\Phi] = \Phi^4$ are:

\begin{enumerate}
  \item $B = 1, g = 1$
  \item $B = 2, g = 0$
\end{enumerate}

Since $E$ is diagonal, the matrix index is conserved along each strand of the ribbon graph. We have to distinguish between internal faces (with constant matrix index) and broken faces which constitute the boundary components. Such a boundary face is characterised by $N_\beta \geq 1$ external double lines to which we attach the source matrices $J$. Conservation of the matrix index along each strand implies that the right index of $J_{ab}$ coincides with the left index of another $J_{bc}$, or of the same $J_{bb}$. Accordingly, the $\beta^{th}$ boundary component carries a cycle $J_{p_1 \ldots p_{N_\beta}} = \prod_{j=1}^{N_\beta} J_{p_j p_{j+1}}$ of $N_\beta$ external source matrices, with $N_\beta + 1 \equiv 1$. 

Being interested in a non-perturbative solution, we will not expand the partition function into ribbon graphs. But we keep the topological information and expand \( \log Z \) according to the cycle structure:

\[
\log \frac{Z[J]}{Z[0]} = \sum_{B=1}^{\infty} \sum_{1 \leq N_1 \leq \cdots \leq N_B} \frac{V^{2-B}}{S_{N_1 \ldots N_B}} \sum_{p_1^\beta \ldots p_{N_\beta}^\beta} G_{[p_1^1 \ldots p_{N_1}^1 \ldots |p_1^\beta \ldots p_{N_\beta}^\beta]} \prod_{\beta=1}^{N_\beta} \frac{J^{N_\beta}}{N_\beta}. 
\]

(6)

The symmetry factor \( S_{N_1 \ldots N_B} \) is obtained as follows: If \( \nu_i \) of the \( B \) numbers \( N_\beta \) in a given tuple \( (N_1, \ldots, N_B) \) are equal to \( i \), then \( S_{N_1 \ldots N_B} = \prod_{\beta=1}^{N_\beta} \nu_i! \).

Next we turn the Ward identity (5) for injective

\[
\sum_{n \in I} \frac{\partial^2 Z[J]}{\partial J_n \partial J_{np}} = \delta_{ap} \left\{ V^2 \sum_{(K)} \frac{J_{P_1} \cdots J_{P_K}}{S_{(K)}} \left( \sum_{n \in I} \frac{G_{[a|n|P_1| \ldots |P_K|}}{V^{|K|+1}} \right) + \sum_{r \geq 1} \sum_{q_1, \ldots, q_r \in I} \frac{G_{[q_1 q_1 \ldots q_r|P_1| \ldots |P_K|} J_{q_1 \ldots q_r}}{V^{|K|+1}} \right. \\
+ V^4 \sum_{(K),(K')} \frac{J_{P_1} \cdots J_{P_K} J_{Q_{K_1}} \cdots J_{Q_{K'}}}{S_{(K)} S_{(K')}} \frac{G_{[a|q_1| \ldots |Q_{K_1}|}}}{V^{|K|+1}} \frac{G_{[a|q_1| \ldots |Q_{K'|}}} V^{|K'|+1}}{V^{|K'|+1}} \right\} Z[J] \\
+ \frac{V}{E_p - E_d} \sum_{n \in I} \left( J_{pn} \frac{\partial Z[J]}{\partial J_n} - J_{na} \frac{\partial Z[J]}{\partial J_{np}} \right) .
\]

(7)

Proof. We identify the following four sources of a singular contribution \( \sim \delta_{ap} \):

1. \( \sum_{n} \frac{\partial^2}{\partial J_n \partial J_{np}} \sum_{q_1, q_2, \ldots} G_{[q_1 q_2 | \ldots |q_1 q_2]} \left( \frac{J_{q_1 q_2}}{2} \right) \prod J \)

2. \( \sum_{n} \frac{\partial^2}{\partial J_n \partial J_{np}} \sum_{q_1, q_2, \ldots} G_{[q_1 | \ldots |q_2]} \left( \frac{J_{q_1 q_2}}{1} \right) \left( \frac{J_{q_2}}{1} \right) \prod J \)

3. \( \sum_{n} \frac{\partial}{\partial J_n} \frac{\partial}{\partial J_{np}} \sum_{q_0, q_1, \ldots, q_{r+1}, \ldots} G_{q_0 q_1 \ldots q_r q_{r+1} | \ldots |q_0 q_1 q_2 \ldots q_{r+1}} \left( \frac{J_{q_0 q_1 q_2 \ldots q_r q_{r+1}}}{r + 2} \right) \prod J \)

\( = \sum_{n} \frac{\partial}{\partial J_n} \sum_{q_1, q_2, \ldots} G_{[p q_1 q_2 | \ldots |q_1 q_2]} \left( \frac{J_{p q_1 q_2}}{1} \right) \prod J \)

4. \( \sum_{n} \frac{\partial}{\partial J_n} \sum_{q_1, q_2, \ldots} G_{[q_1 q_2 | \ldots |q_1 q_2]} \left( \frac{J_{q_1 q_2}}{1} \right) \prod J \)
4. \[ \sum_n \frac{\partial^2}{\partial J_{an} \partial J_{np}} \left[ \sum_{q_1} G_{n|q_1} (J_{q_1}) \prod J \right] \left[ \sum_{q_2} G_{n|q_2} (J_{q_2}) \prod J \right] \]

All other types of derivatives, collected into \( \left( \sum_{n \in I} \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right)_{\text{reg}} \), persist for \( a \neq p \). For \( p \neq a \) we clearly have

\[ \left( \sum_{n \in I} \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right)_{\text{reg}} = \sum_{n \in I} \left. \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right|_{a \neq p} = \frac{V}{E_p - E_a} \left( J_{pa} \frac{\partial Z}{\partial J_{an}} - J_{na} \frac{\partial Z}{\partial J_{np}} \right), \] (8)

where the last equality is the Ward identity \( \Box \), divided by \( \frac{E_p - E_a}{V} \neq 0 \). By a continuity argument, the rightmost term in \( \Box \) must agree with \( \left( \sum_{n \in I} \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right)_{\text{reg}} \) also in the limit \( p \to a \), and this finishes the proof.

2.2 Schwinger-Dyson equations

We can write the action as \( S = \frac{V}{2} \sum_{a,b} (E_a + E_b) \Phi_{ab} \Phi_{ba} + V S_{\text{int}}[\Phi] \), where \( E_a \) are the eigenvalues of \( E \). Functional integration yields, up to an irrelevant constant,

\[ Z[J] = e^{-V S_{\text{int}}[\frac{\Phi}{\sqrt{V} J}]} e^{\frac{V}{2} \langle J, J \rangle_E} \]

\[ \langle J, J \rangle_E := \sum_{m,n \in I} \frac{J_{mn} J_{nm}}{E_m + E_n}. \] (9)

Instead of a perturbative expansion of \( e^{-V S_{\text{int}}[\frac{\Phi}{\sqrt{V} J}]} \), we apply those \( J \)-derivatives to \( \Box \) which give rise to a correlation function \( G_- \) on the lhs. On the rhs of \( \Box \), these external derivatives combine with internal derivatives from \( S_{\text{int}}[\frac{\Phi}{\sqrt{V} J}] \) to certain identities for \( G_- \). These Schwinger-Dyson equations are often of little use because they express an \( N \)-point function in terms of \((N+2)\)-point functions.

In the field-theoretical matrix models under consideration, the Ward identity \( \Box \) lets this tower of Schwinger-Dyson equations collapse. To see this we consider the 2-point function \( G_{[ab]} \) for \( a \neq b \). According to \( \Box \), \( G_{[ab]} \) is obtained by deriving \( \Box \) with respect to \( J_{ba} \) and \( J_{ab} \):

\[ G_{[ab]} = \frac{1}{V Z[0]} \left. \frac{\partial^2 Z[J]}{\partial J_{ba} \partial J_{ab}} \right|_{J=0} \] (disconnected part of \( Z \) does not contribute for \( a \neq b \))

\[ = \frac{1}{V Z[0]} \left\{ \frac{\partial}{\partial J_{ba}} e^{-V S_{\text{int}}[\frac{\Phi}{\sqrt{V} J}]} e^{\frac{V}{2} \langle J, J \rangle_E} \right\}_{J=0} \]

\[ = \frac{1}{(E_a + E_b) Z[0]} \left\{ \frac{\partial}{\partial J_{ba}} e^{-V S_{\text{int}}[\frac{\Phi}{\sqrt{V} J}]} J_{ba} e^{\frac{V}{2} \langle J, J \rangle_E} \right\}_{J=0} \]

\[ = \frac{1}{E_a + E_b} + \frac{1}{(E_a + E_b) Z[0]} \left\{ \left( \Phi_{ab} \frac{\partial}{\partial \Phi_{ab}} \right) \left[ \frac{\partial}{\partial \Phi_{ab}} \right] \right\} Z[J]_{J=0}. \] (10)

Now observe that \( \frac{\partial}{\partial \Phi_{ab}} (-V S_{\text{int}}) \) contains, for any \( P[\Phi] \), the derivative \( \sum_n \frac{\partial^2}{\partial J_{an} \partial J_{np}} \) which we know from \( \Box \). In case of the quartic matrix model \( P[\Phi] = \frac{1}{4} \Phi^4 \) we
have $\frac{\partial (-V S_{int})}{\partial \Phi_{ab}} = -\lambda V \sum_{n,p \in I} \Phi_{bp} \Phi_{pn} \Phi_{na}$, hence

$$\left( \Phi_{ab} \frac{\partial (-V S_{int})}{\partial \Phi_{ab}} \right) \left[ \frac{\partial }{\partial J} \right] = -\frac{\lambda}{V^3} \sum_{p,n \in I} \frac{\partial^2}{\partial J_{pb} \partial J_{na}} \frac{\partial^2}{\partial J_{an} \partial J_{np}},$$

and the Schwinger-Dyson equation \((10)\) for $G_{(ab)}$ becomes with \((7)\)

$$G_{(ab)} = \frac{1}{E_a + E_b} - \frac{\lambda}{V^3(E_a + E_b)Z[0]} \sum_{p \in I} \frac{\partial^2}{\partial J_{pb} \partial J_{na}} \frac{\partial^2}{\partial J_{an} \partial J_{np}} \bigg|_{J=0}
$$

$$= \frac{1}{E_a + E_b} - \frac{\lambda}{V(E_a + E_b)Z[0]} \frac{\partial}{\partial J} \left\{ \sum_{n \in I} G_{[an]} \frac{V}{V^2} - \frac{G_{[an]q}}{V^3} \sum_{n,q,r \in I} G_{[an]qr} J_{qr} J_{rr} \right.$$

$$+ \sum_{q,r \in I} G_{[qq]} J_{qr} J_{rr} + \frac{V^2}{V^2} G_{[q]q} \frac{G_{[q]} J_{rr}}{1} \bigg|_{J=0}
$$

$$- \frac{\lambda}{V^2(E_a + E_b)Z[0]} \sum_{p \in I} \left( \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{na}} + \frac{\partial^2 Z[J]}{\partial J_{pb} \partial J_{na}} - \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right) \bigg|_{J=0}.$$

Taking $\frac{\partial^2 Z[J]}{\partial J_{pb} \partial J_{na}} = (V G_{[pb]} + \delta_{pb} G_{[p]b})Z[0] + \mathcal{O}(J)$ and $\frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} = 0$ for $a \neq b$ into account, we have proved:

**Proposition 3** The 2-point function of a quartic matrix model with action $S = V \text{tr}(E \Phi^2 + \frac{1}{4} \Phi^4)$ satisfies for injective $m \mapsto E_m$ the Schwinger-Dyson equation

$$G_{[ab]} = \frac{1}{E_a + E_b} - \frac{\lambda}{V^2(E_a + E_b)} G_{[an]} \frac{V}{V^2} - \frac{G_{[an]q}}{V^3} \sum_{n,q,r \in I} G_{[an]qr} J_{qr} J_{rr}$$

$$+ \frac{V^2}{V^2} G_{[q]q} \frac{G_{[q]} J_{rr}}{1}$$

$$- \frac{\lambda}{V^2(E_a + E_b)} \sum_{p \in I} \left( \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{na}} + \frac{\partial^2 Z[J]}{\partial J_{pb} \partial J_{na}} - \frac{\partial^2 Z[J]}{\partial J_{an} \partial J_{np}} \right) \bigg|_{J=0}.$$

It can be checked [GW12b] that in a genus expansion $G_{\ldots} = \sum_{g=0}^{\infty} V^{-2g} G^{(g)}$ (which is probably not convergent but Borel summable), precisely the line \((12a)\) preserves the genus, the lines \((12b)\) increase $g \mapsto g + 1$ and the line \((12c)\) increases $g \mapsto g + 2$. In particular, in a scaling limit $V \to \infty$ with $\frac{1}{V} \sum_{p \in I}$ finite, the exact
Schwinger-Dyson equation for $G_{[ab]}$ coincides with its restriction (12a) to the planar sector $g = 0$, a closed non-linear equation for $G^{(0)}_{[ab]}$ alone:

$$G^{(0)}_{[ab]} = \frac{1}{E_a + E_b} - \frac{\lambda}{E_a + E_b} \frac{1}{V} \sum_{p \in I} \left( G^{(0)}_{[ap]} G^{(0)}_{[ab]} - \frac{G^{(0)}_{[pb]} - G^{(0)}_{[ab]}}{E_p - E_a} \right). \quad (13)$$

We have derived in 2007/08 this self-consistency equation for the Moyal model by the graphical method proposed by [DGMR07]. In this form, (13) is meaningless because $\sum_{p \in I}$ diverges. In 2009 we solved the renormalisation problem, namely the renormalisation of infinitely many Feynman graphs at once [GW09]. This renormalisation increases the non-linearity. In [GW09] we have solved (13) perturbatively to $O(\lambda^3)$. After several years of setbacks with the non-perturbative solution, a breakthrough came in 2012: The equation (13) can be turned into an equation which is linear in the difference $G^{(0)}_{[ab]} - G^{(0)}_{[ab]}$ to the boundary and non-linear only in $G^{(0)}_{[ab]}$!

A similar calculation gives the Schwinger-Dyson equation for higher $N$-point functions:

$$G^{(0)}_{[a_1 \ldots a_{N-1}]} = -\frac{\lambda}{V} \sum_{p \in I} \left( G_{[ap]} G_{[a_1 \ldots a_{N-1}]} - \frac{G_{[pb]} - G_{[ab]} |}{E_p - E_a} \right) \quad (14a)$$

$$-\frac{\lambda}{V^2 (E_a + E_b)} \left( G_{[a|a]} G_{[a_1 \ldots a_{N-1}]} + \sum_{k=1}^{N-1} G_{[b_1 \ldots b_k a b_k \ldots b_{N-1}] a} \right) \quad (14b)$$

$$-\frac{\lambda}{V^4 (E_a + E_b)} G_{[a|a]} |a_1 \ldots a_{N-1}|. \quad (14c)$$

Again, the first lines (14a) preserve the genus, whereas $g \mapsto g+1$ in (14b) and $g \mapsto g+2$ in (14c). The planar sector $G^{(0)}_{[a_1 \ldots a_{N-1}]}$, exact for $V \to \infty$ with $\frac{1}{V} \sum_{p \in I}$ finite, is a linear inhomogeneous equation with inductively known parameters.

It turns out that a real theory with $\Phi = \Phi^*$ admits a short-cut which directly gives the higher $N$-point functions without any index summation. Since the equations for $G_{[aa]}$ are real and $J_{ab} = J_{ba}$, the reality $\mathcal{Z} = \overline{\mathcal{Z}}$ implies (in addition to invariance under cyclic permutations) invariance under orientation reversal

$$G_{[p_1 | \ldots | p_{k_1-1} | - | p_{k_2} | p_{k_1} | \ldots | p_{N-1} |] = G_{[p_1 | \ldots | p_{k_1-1} | - | p_{k_2} | p_{k_1} | \ldots | p_{N-1} |].} \quad (15)$$
The last line of (16) increases the genus and is absent in $B = 2$ boundary components. 

Write down the same equation but with $\lambda$ to together with a corresponding rescaling $\lambda \mapsto Z \lambda$ leaves the algebraic equations (16) as well as (65) and (66) invariant:

**Proposition 4**

$$G_{[b_0 b_1 \ldots b_{N-1}]} = \left( -\lambda \right) \sum_{l=1}^{N+2} \frac{G_{[b_0 b_1 \ldots b_{2l-1}]}G_{[b_2 b_3 \ldots b_{N-1}]} - G_{[b_2 b_3 \ldots b_{2l-1}]}G_{[b_0 b_1 \ldots b_{2l-1}]} }{ (E_{b_0} - E_{b_2}) (E_{b_1} - E_{b_{N-1}}) }$$

$$+ \left( -\lambda \right) \frac{1}{V^2} \sum_{k=1}^{N-1} \frac{G_{[b_0 b_1 \ldots b_k b_{k+1} \ldots b_{N-1}]} - G_{[b_k b_1 \ldots b_{k-1} b_{k+1} \ldots b_{N-1}]} }{ (E_{b_0} - E_{b_k}) (E_{b_1} - E_{b_{N-1}}) } .$$

(16)

The last line of (16) increases the genus and is absent in $G_{[b_0 b_1 \ldots b_{N-1}]}^{(0)}$. Instead of giving the general proof, let us look at the case $N = 4$. Then (16), multiplied by $E_a - E_{b_1}$, reads

$$(E_a - E_{b_1}) G_{[abcd]}$$

$$= \left( -\lambda \right) \left( \frac{1}{V} \sum_{p \in I} \left( G_{[ap]} G_{[abcd]} - \frac{G_{[pbcd]} - G_{[abcd]}}{E_p - E_a} \right) - G_{[bc]} G_{[da]} - G_{[cd]} \right)$$

$$- \frac{\lambda}{V^2} \left( G_{[a|a]} G_{[abcd]} + G_{[bc|da]} + G_{[bd|ca]} + G_{[cd|ab]} + \frac{1}{V} \sum_{p \in I} G_{[ap|abcd]} \right)$$

$$- \frac{G_{[bc|da]} - G_{[bd|ca]} - G_{[bd|ca]} - G_{[cd|ab]} - G_{[cd|ab]} - G_{[a|a|abcd]} }{E_b - E_a}$$

$$- \frac{G_{[bd|ca]} - G_{[bd|ca]} - G_{[cd|ab]} - G_{[a|a|abcd]} }{E_c - E_a}$$

$$- \frac{G_{[cd|ab]} - G_{[a|a|abcd]} }{E_d - E_a} .$$

(17)

Write down the same equation but with $b \leftrightarrow d$, and take the difference between these equations. Then most terms cancel because by (15) we have the equalities $G_{[abcd]} = G_{[adcb]}$, $G_{[pbcd]} = G_{[pdcb]}$, $G_{[bc|da]} = G_{[dc|ba]}$, $G_{[bd|ca]} = G_{[da|cb]}$, $G_{[bd|ca]} = G_{[dca|ba]}$, $G_{[bc|da]} = G_{[dca|ba]}$, $G_{[bd|ca]} = G_{[da|cb]}$, $G_{[bd|ca]} = G_{[dca|ba]}$, $G_{[bd|ca]} = G_{[adcb]}$, $G_{[bc|da]} = G_{[a|a|abcd]}$ and $G_{[a|a|abcd]} = G_{[a|a|abcd]}$. Altogether, the difference (17) - (17) vanishes upon cancellation

$$(E_d - E_{b_1}) G_{[abcd]} = \left( -\lambda \right) \frac{G_{[ab]} G_{[cd]} - G_{[ad]} G_{[cb]} }{E_c - E_a}$$

$$- \frac{\lambda}{V^2} \left( G_{[b|da]} - G_{[a|a|abcd]} \right) E_b - E_a$$

and this is (16) for $N = 4$.

For completeness, we list in the appendix the Schwinger-Dyson equation for $B = 2$ boundary components.

We make the following key observation: An affine transformation $E \mapsto Z E + C$ together with a corresponding rescaling $\lambda \mapsto Z^2 \lambda$ leaves the algebraic equations (16) as well as (65) and (66) invariant:
Theorem 5 Given a real quartic matrix model with $S = V \tr (E \Phi^2 + \frac{\lambda}{4} \Phi^4)$ and $m \mapsto E_m$ injective, which determines the set $G_{|p_1^1 \ldots p_N^{N_1}| \ldots |p_B^1 \ldots p_B^{N_B}|}$ of $(N_1 + \ldots + N_B)$-point functions. Assume that the basic functions with all $N_i \leq 2$ are turned finite by $E_a \mapsto Z(E_a + \frac{\mu^2}{2} - \frac{\mu_{bare}^2}{2})$ and $\lambda \mapsto Z^2 \lambda$. Then all functions with one $N_i \geq 3$ are finite without further need of a renormalisation of $\lambda$, i.e. all renormalisable quartic matrix models have vanishing $\beta$-function.

The theorem tells us that vanishing of the $\beta$-function for the self-dual $\Phi^4$-model on Moyal space (proved in [DGMR07] to all orders in perturbation theory) is generic to all quartic matrix models, and the result even holds non-perturbatively!

The universal recursion formula (16) computes the planar $N$-point function $G_{|b_0 \ldots b_{N-1}|}$ at $B = 1$ as a sum of fractions with products of 2-point functions in the numerator and products of differences of eigenvalues of $E$ in the denominator. This structure admits an interesting graphical interpretation. We draw the indices $b_0, \ldots, b_{N-1}$ in cyclic order on the circle $S^1$ and represent a factor $G_{b_i b_j}$ as a chord connecting $b_i$ with $b_j$ and a factor $\frac{1}{E_{b_i} - E_{b_j}}$ as an arrow from $b_i$ to $b_j$:

$$G_{b_0 b_1 b_2 b_3}^{(0)} = (-\lambda) \frac{G_{b_0 b_1} G_{b_2 b_3} - G_{b_0 b_2} G_{b_1 b_3}}{(E_{b_0} - E_{b_1})(E_{b_1} - E_{b_3})} = (-\lambda) \left\{ \begin{array}{c}
\includegraphics[width=0.3\textwidth]{fig16}
\end{array} \right\},$$

$$G_{b_0 \ldots b_5} = (-\lambda)^2 \left\{ \begin{array}{c}
\includegraphics[width=0.6\textwidth]{fig18}
\end{array} \right\}.$$  

The chords form the non-crossing chord diagrams counted by the Catalan number $C_N = \frac{N!}{\left( \frac{N}{2} + 1 \right)! \frac{N}{2}!}$. The arrows form two disjoint trees, one connecting the even vertices and one connecting the odd vertices. By rational fraction expansion it is possible to achieve that each tree intersects the chord only in the vertices.
The assignment of trees to a given chord diagram is, in general, not unique. A canonical choice is not known to us.

2.3 Digression: Quantum gravity in two dimensions

Two-dimensional quantum gravity (see [DGZ95, ADJ97] for reviews) can be interpreted as the enumeration of random triangulations of surfaces. Its asymptotic behaviour is captured by the matrix model partition function

\[ Z = \int \mathcal{D}[\Phi] \exp \left( - \mathcal{N} \sum_n t_n \text{tr}(\Phi^n) \right), \tag{19} \]

where the integral is over \((\mathcal{N} \times \mathcal{N})\)-Hermitean matrices \(\Phi\) and the \(t_n\) are scalar coefficients. In the limit \(\mathcal{N} \to \infty\), this series in \((t_n)\) is evaluated in terms of the \(\tau\)-function for the Korteweg-de Vries (KdV) hierarchy. There is another approach to topological gravity in which the partition function is a series in \((t_n)\) with coefficients given by intersection numbers of complex curves. Witten conjectured [Wit91] that the partition functions of the two approaches coincide. This conjecture was proved by Kontsevich [Kon82] who achieved the computation of the intersection numbers in terms of weighted sums over ribbon graphs (fat Feynman graphs), which he proved to be generated from the Airy function matrix model (Kontsevich model)

\[ Z[E] = \frac{\int \mathcal{D}[\Phi] \exp \left( - \frac{1}{2} \text{tr}(E\Phi^2) + \frac{1}{6} \text{tr}(\Phi^3) \right)}{\int \mathcal{D}[\Phi] \exp \left( - \frac{1}{2} \text{tr}(E\Phi^2) \right)}. \tag{20} \]

The external matrix \(E = E^* > 0\) is related by \(t_n = (2n-1)!! \text{tr}(E^{-(2n-1)})\) to the series \((t_n)\). The limit \(\mathcal{N} \to \infty\) of \(Z[E]\) gives the KdV evolution equation, thus proving Witten’s conjecture.

We have proved that also the quartic matrix model

\[ Z[E, J, \lambda] = \frac{\int \mathcal{D}[\Phi] \exp \left( - \text{tr}(E\Phi^2) + \text{tr}(J\Phi) - \frac{\lambda}{4} \text{tr}(\Phi^4) \right)}{\int \mathcal{D}[\Phi] \exp \left( - \text{tr}(E\Phi^2) - \frac{\lambda}{4} \text{tr}(\Phi^4) \right)}, \tag{21} \]

is in the large-\(\mathcal{N}\) limit exactly solvable in terms of the solution of a non-linear equation [13]. Any triangulation can be subdivided into a quadrangulation

![Diagram: Quadrangulation](and vice versa). From Witten’s uniqueness argument [Wit91], 2D quantum gravity should have equivalent descriptions as cubic [20] and quartic [21] matrix model. Understanding the precise relation between [20] and [21] would be of high interest:
1. In contrast to (21), the cubic action (20) lacks manifest positivity due to its purely imaginary coupling constant.

2. A quartic action admits a Hubbard-Stratonovich transform which is the key ingredient of a new approach to constructive quantum field theory [Riv07b] that avoids the cluster expansion.

3. Conversely, the integrability of (20) might provide valuable information about the solution of the self-consistency equation (13).

Coloured tensor models (see [GP12, Riv13] for recent reviews) extend these methods to quantum gravity in $D \geq 3$. They became a very active domain of research after understanding [Gur10] of the analogue of the large-$N$ behaviour of matrix models [tHo74]. They have Schwinger-Dyson equations (see e.g. [Bon12]) and action of the $U(\infty)$ group. A first promising result is the recent derivation of closed equations for the 2-point functions of rank 3 and 4 tensorial group field theory [Sam14].

3 $\Phi_4^3$-theory on Moyal space as a fixed point problem

3.1 Preliminaries

Taking the renormalisation group [WK74] serious, we would expect that General Relativity, because not renormalisable, is irrelevant and hence scaled away. The existence of gravity thus tells us that the scaling must stop at some length scale, and from the weakness of the gravitational coupling constant one deduces the value of that scale: the Planck length $10^{-35}$ m. There, the geometry of nature is expected to differ from the familiar structure of a differentiable manifold. One of many candidates for Planck scale physics is noncommutative geometry [Con94], a vast reformulation of geometry and topology in the language of operator algebras. The focus is shifted from manifolds to generalisations of the algebra of functions. This concept proved very successful in understanding the geometry of the Standard Model of particle physics as Riemannian geometry of a space which is the product of a manifold with a discrete space [Con96, CC96].

A large class of examples of noncommutative geometries comes from deformations of the algebra of functions on manifolds. Schwartz functions on Euclidean space $\mathbb{R}^4$ admit an $\mathbb{R}^4$-group action by translation. As shown by Rieffel [Rie93], this group action induces a noncommutative associative product on the space of Schwartz functions, the Moyal product:

$$(f \star g)(x) = \int_{\mathbb{R}^4 \times \mathbb{R}^4} \frac{dy \, dk}{(2\pi)^4} f(x+\frac{i}{2} \Theta kj) \, g(x+y) \, e^{ij(k,y)} , \quad \Theta = -\Theta^t \in M_4(\mathbb{R}). \quad (22)$$

Whether or not the Moyal space ($\mathbb{R}^4, \star$) is relevant for Planck scale physics is pure speculation (although a refinement can be justified by uncertainty relations for position operators [DFR95]). In any case the Moyal space is a nice toy
model on which it is easy to formulate and to study (quantum) field theories. To formulate a Euclidean quantum field theory on Moyal space it is, at first sight, enough to replace in the action of a usual field theory the pointwise product of functions by the $\star$-product. The simplest example is the $\phi^4$-model with action
\begin{equation}
S[\phi] = \int_{\mathbb{R}^4} dx \left( \frac{1}{2} \phi \ast (-\Delta + \mu^2)\phi + \frac{\lambda}{4} \phi \ast \phi \ast \phi \ast \phi \right)(x) . \tag{23}
\end{equation}

The resulting Feynman rules [Fil96] lead to situations where a multiple insertion of non-planar subgraphs gives rise to divergences of arbitrarily high degree (ultraviolet/infrared mixing [MVS00]). See [CR00] for a thorough investigation of this problem. Relativistic quantum field theories on noncommutative Minkowski space are much more difficult [BDFP02]. Here the UV/IR-mixing problem occurs in different types of graphs [Bah10].

The Moyal algebra $(\mathcal{S}(\mathbb{R}^4), \ast)$ has a matrix basis [GV88, VG88, GGISV03]
\begin{equation}
\phi(x) = \sum_{m,n \in \mathbb{N}^2} \Phi_{mn} f_{mn}(x), \quad f_{mn}(x) = f_{m_1n_1}(x^0,x^1)f_{m_2n_2}(x^3,x^4) ,
\end{equation}
\begin{equation}
f_{mn}(y^0,y^1) = 2(-1)^m \sqrt{m!} (\sqrt{2} \theta)^{n-m} L_n^m \left( \frac{|y|^2}{\theta} \right) e^{-\frac{|y|^2}{\theta}} , \tag{24}
\end{equation}
where $L_n^m$ are Laguerre polynomials, $y \equiv y^0+iy^1$ and $\underline{m} = (m_1,m_2)$. Without loss of generality we assume the only non-vanishing components of $\Theta$ to be $\Theta_{12} = -\Theta_{21} = \Theta_{34} = -\Theta_{43}$. The functions $f_{mn}$ satisfy
\begin{equation}
(f_{kl} \ast f_{mn})(x) = \delta_{kl} f_{mn}(x), \quad \int_{\mathbb{R}^4} dx f_{mn}(x) = (2\pi \theta)^2 \delta_{mn} .
\end{equation}

Therefore, the $\phi^4$-interaction in (23) becomes a matrix product (we write $\phi$ for a function and $\Phi$ for a matrix):
\begin{equation}
S[\phi] = (2\pi \theta)^2 \sum_{k,l,m,n \in \mathbb{N}^2} \left( \frac{1}{2} \Phi_{kl} (\Delta_{kl,mn} + \mu^2 \delta_{kn} \delta_{lm}) \Phi_{mn} + \frac{\lambda}{4} \Phi_{kl} \Phi_{lm} \Phi_{mn} \Phi_{nk} \right) . \tag{25}
\end{equation}

The matrix kernel $\Delta_{kl,mn}$ of the Laplacian $(-\Delta)$, viewed as map from $\mathbb{N}^4$ to $\mathbb{N}^4$, consists of a local interaction plus nearest neighbour interaction.

In [GW05b] we studied the renormalisation group flow of the $\phi^4$-model in matrix representation (using a power-counting theorem [GW05a] for matrix models with kernel $\Delta_{kl,mn}$). We noticed that the marginal parts of the local term and of the nearest neighbour term in $\Delta_{kl,mn}$ have different flows. To absorb these different flows a 4th relevant/marginal operator in the action functional is necessary. This operator corresponds to a harmonic oscillator potential:
\begin{equation}
S[\phi] = 64\pi^2 \int d^4x \left( Z^2 \phi \ast (-\Delta + \Omega^2 (2\Theta^{-1} x)^2 + \mu_{bare}^2) \phi + \frac{\lambda Z^2}{4} \phi \ast \phi \ast \phi \ast \phi \right)(x) . \tag{26}
\end{equation}
We proved in [GW05b] that the corresponding Euclidean quantum field theory is renormalisable to all orders in perturbation theory. This result was reestablished by various methods, see [Riv07a] for a review.

Presence of the harmonic oscillator term $\Omega \neq 0$ breaks translation invariance. Conversely, this term achieves covariance under Langmann-Szabo duality transformation [LS02] which consists in exchanging $x \leftrightarrow p$ and $\phi(x) \leftrightarrow \phi(p)$ followed by Fourier transform back to the original variables. Remarkably, this transformation leaves $\int dx \phi^* \phi^* \phi^* \phi^* \phi$ invariant, and it exchanges $\int dx \phi(-\Delta)\phi$ with $\int dx \phi|2\Theta^{-1}x|^2 \phi$. Presence of the oscillator term gives rise to an interesting spectral noncommutative geometry [GW13a] (see also [GW12a]) which is conceptually simpler than the isospectral deformation [GGISV03] of $\mathbb{R}^4$. Most importantly, the oscillator term cures the Landau ghost problem [LAK54a, LAK54b, LAK54c] of usual $\phi^4_4$-theory: We have discovered in [GW04, GW05c] that the one-loop renormalisation group flows of $\Omega$ and $\lambda$ influence each other in such a way that the running coupling constant $\lambda(\Lambda)$ remains finite at any scale $\Lambda$. Even more, at the self-duality point $\Omega = 1$ the $\beta$-function of the $\lambda \Phi^4_4$-coupling vanishes to all orders in perturbation theory [DGMR07]. This result was obtained by an ingenious combination of Ward identities and Schwinger-Dyson equations (see [DR07] for an explicit three-loop calculation). In [GW12b] we have generalised the method of Disertori-Gurau-Magnen-Rivasseau [DGMR07] to the whole class of quartic matrix models (reviewed in sec. 2). Vanishing of the $\beta$-function is often connected with integrability, and together with the absent Landau ghost problem a non-perturbatively constructed $\phi^4_4$-model on Moyal space came into reach. The first milestone was the derivation of the self-consistency equation (13) and the understanding of its renormalisation in [GW09]. It took us several years to fully understand this equation, and it is only recently that we finished the solution/construction of the Moyal space $\phi^4_4$-model [GW12b]. In the sequel we review this construction.

3.2 Renormalisation and integral representation

At the self-duality point $\Omega = 1$, the matrix kernel $\Delta_{k,l|mn}^{\Omega=1}$ of the Schrödinger operator $H = -\Delta + ||2\Theta^{-1}x||^2$ becomes purely local and turns the action (26) in matrix basis (24) into a (field-theoretical) quartic matrix model with action

$$S[\Phi] = V \left( \sum_{m,n \in \mathbb{N}^2_N} E_m \Phi_{mn} \Phi_{mn} + \frac{Z^2 \lambda}{4} \sum_{m,n,k,l \in \mathbb{N}^2_N} \Phi_{mn} \Phi_{nk} \Phi_{kl} \Phi_{lm} \right),$$

$$E_m = Z \left( \frac{|m|}{\sqrt{V}} + \frac{\mu_{bare}}{2} \right), \quad |m| := m_1 + m_2 \leq N, \quad V = \left( \frac{\theta}{4} \right)^2.$$
Our general results on quartic matrix models imply that the planar 2-point function $G^{(0)}_{ab}$ satisfies the self-consistency equation (13),

$$G^{(0)}_{ab} = \frac{1}{E_a + E_b} - \frac{Z^2 \lambda}{V} \sum_{p \in N_N^2} \left( \frac{G^{(0)}_{ab} G^{(0)}_{ap}}{E_p - E_a} - \frac{G^{(0)}_{ab} G^{(0)}_{bp}}{E_p - E_b} \right). \quad (28)$$

We have introduced a cut-off $N_N^2$ in the matrix size; the index sum diverges for $N_N^2 \rightarrow N^2$. As usual, the renormalisation strategy consists in adjusting $Z, \mu_{\text{bare}}$ in such a way that the limit $N_N^2 \rightarrow N^2$ exists. This will be achieved by normalisation conditions for the 1PI function $\Gamma_{ab}$ defined by $G^{(0)}_{ab} =: (H_{ab} - \Gamma_{ab})^{-1}$, where $H_{ab} := E_a + E_b$. We express (28) in terms of $\Gamma_{ab}$,

$$\Gamma_{ab} = -\frac{\lambda Z^2}{V} \sum_{p \in N_N^2} \left( \frac{1}{H_{ap} - \Gamma_{ap}} + \frac{1}{H_{bp} - \Gamma_{bp}} - \frac{1}{(H_{ap} - \Gamma_{ap})(H_{bp} - \Gamma_{bp})} \right), \quad (29)$$

and write $\Gamma_{ab}$ as first-order Taylor formula with remainder $\Gamma_{ab}^{\text{ren}}$,

$$\Gamma_{ab} = Z \mu_{\text{bare}}^2 - \mu^2 + \frac{(Z - 1)}{\sqrt{V}} (|a| + |b|) + \Gamma_{ab}^{\text{ren}}, \quad \Gamma_{ab}^{\text{ren}}_{oo} = 0, \quad (\partial \Gamma_{ab}^{\text{ren}}_{oo})_{ab} = 0.$$

Equation (29) for $\Gamma_{ab} \left[ \Gamma_{ab}^{\text{ren}}, \mu_{\text{bare}}^2, Z \right]$ together with $\Gamma_{oo}^{\text{ren}} = 0$ and $(\partial \Gamma_{ab}^{\text{ren}})^{oo}_{ab}$ constitute three equations to determine the three functions $\Gamma_{ab}^{\text{ren}}, \mu_{\text{bare}}^2, Z$. Eliminating $\mu_{\text{bare}}^2, Z$ thus gives rise to a closed equation for the renormalised function $\Gamma_{ab}^{\text{ren}}$ alone. For this elimination it is important to note that the equations for $\Gamma_{ab}^{\text{ren}}, \mu_{\text{bare}}, Z$ depend on $a, b$ only via the norms $|a|, |b|$ which parametrise the spectrum of $E$. Therefore, $\Gamma_{ab}$ is actually a function only of $|a|, |b|$, and consequently the index sum reduces to $\sum_{p \in N_N^2} f(|p|) = \sum_{|p| = 0}^N (|p| + 1) f(|p|)$.

We study a particular scaling limit in which matrix size $N$ and volume $V$ are simultaneously sent to $\infty$ such that the ratio $\frac{N}{\sqrt{V} \mu^4} = \Lambda^2 (1 + Y)$ is kept fixed. Note that $V = \left(\frac{N}{\Lambda^2}\right)^2 \rightarrow \infty$ is a limit of extreme noncommutativity! The new parameter $(1 + Y)$ corresponds to a finite wavefunction renormalisation, identified later to decouple our equations. The parameter $\Lambda^2$ represents an ultraviolet cut-off which is sent to $\Lambda \rightarrow \infty$ in the very end (continuum limit). In the scaling limit, functions of $\frac{|p|}{\sqrt{V}} =: \frac{\mu^2 (1 + Y)}{\Lambda^2}$ converge to functions of `continuous matrix indices’ $p \in [0, \Lambda^2]$. In the same way, $\Gamma_{ab}^{\text{ren}}$ converges to a function $\mu^2 \Gamma_{ab}$ with $a, b \in [0, \Lambda^2]$, and the discrete sum converges to a Riemann integral

$$\frac{1}{V} \sum_{|p| = 0}^N (|p| + 1) f \left( \frac{|p|}{\Lambda} \right) \rightarrow \mu^4 (1 + Y)^2 \int_0^{\Lambda^2} p \, dp \, f \left( \mu^2 (1 + Y) p \right).$$

This limit makes the restriction to the planar sector (13) of (12) exact.
After elimination of $\mu^2_{\text{bare}}$, but before elimination of $Z$, our equation for $\Gamma_{ab}$ becomes

\[
(Z - 1)(1 + Y)(a + b) + \Gamma_{ab} = -\lambda (1 + Y) \int_0^{\Lambda^2} dp \left( \frac{Z^2}{(a + p)(1 + Y) + 1 - \Gamma_{ap}} - \frac{Z^2}{p(1 + Y) + 1 - \Gamma_{0p}} \right) - \lambda (1 + Y)^2 \int_0^{\Lambda^2} dp \left( \frac{Z}{(b + p)(1 + Y) + 1 - \Gamma_{pb}} - \frac{Z}{p(1 + Y) + 1 - \Gamma_{p0}} \right) - \frac{(b + p)(1 + Y) \Gamma_{pb} - \Gamma_{ab}}{(1 + Y)(p - a)} \right) + \Gamma_{ab} - \Gamma_{00} \cdot \frac{(1 + Y)(a + b)}{\lambda}.
\] (30)

Applying $\frac{d}{db}|_{a=b=0}$, we get $Z$ in terms of $\Gamma_{ab}$ (and its derivative). Inserted back one gets a highly non-linear integro-differential equation. Fortunately we can reduce the non-linearity by subtracting from (30) the same equation taken at $b = 0$. This subtraction eliminates the second line of (30) containing $Z^2$. In terms of $G_{ab} := \frac{(a + b)(1 + Y) + 1 - \Gamma_{ab}}{-\lambda}$, this difference equation reads

\[
(1 + Y)^{-1} \left( \frac{1}{G_{ab}} - \frac{1}{G_{a0}} \right) = b - \lambda \int_0^{\Lambda^2} dp \frac{G_{pb} - G_{p0}}{p - a} \cdot \frac{G_{p0}}{G_{a0}}.
\] (31)

Differentiation $\frac{d}{db}|_{a=b=0}$ of (31) yields $Z$ in terms of $G_{ab}$ and its derivative. The resulting derivative $G'$ can be avoided by adjusting

\[
Y := -\lambda \lim_{b \to 0} \int_0^{\Lambda^2} dp \frac{G_{pb} - G_{p0}}{b}.
\]

This choice leads to

\[
\frac{Z^{-1}}{(1 + Y)} = 1 - \lambda \int_0^{\Lambda^2} dp G_{p0},
\]

which is a perturbatively divergent integral for $\Lambda \to \infty$. Inserting $Z^{-1}$ and $Y$ back into (31) we end up in a linear integral equation for the difference function $D_{ab} := \frac{2}{\lambda}(G_{ab} - G_{a0})$ to the boundary:

\[
\left( \frac{b}{a} + \frac{1}{aG_{a0}} \right) D_{ab} + G_{a0} = \lambda \int_0^{\Lambda^2} dp \frac{D_{pb} - D_{0b}G_{pb}}{p - a}.
\] (32)

The non-linearity restricts to the boundary function $G_{a0}$ where the second index is put to zero. Assuming $a \mapsto G_{ab}$ Hölder-continuous, we can pass to Cauchy principal values. In terms of the finite Hilbert transform

\[
\mathcal{H}_a^\Lambda[f(\bullet)] := \frac{1}{\pi} \lim_{\epsilon \to 0} \left( \int_0^{a-\epsilon} + \int_{a+\epsilon}^{\Lambda^2} \right) \frac{f(q) dq}{q - a},
\] (33)
the integral equation (32) becomes
\[
\left( \frac{b}{a} + 1 + \frac{\lambda \pi a \mathcal{H}_a^{\lambda} [G_{\bullet 0}]}{a G_{a 0}} \right) D_{ab} - \lambda \pi \mathcal{H}_a^{\lambda} [D_{\bullet b}] = -G_{a 0}. \tag{34}
\]

3.3 The Carleman solution

Equation (34) is a well-known singular integral equation of Carleman type [Car22, Tri57]:

Theorem 6 ([Tri57], transformed from \([-1, 1]\) to \([0, \Lambda^2]\)) The singular linear integral equation
\[
h(a)y(a) - \lambda \pi \mathcal{H}_a^{\lambda}[y] = f(a), \quad a \in ]0, \Lambda^2[,
\]
is for \(h(a)\) continuous on \([0, \Lambda^2]\), Hölder-continuous near \(0, \Lambda^2\), and \(f \in L^p\) for some \(p > 1\) (determined by \(\vartheta(0)\) and \(\vartheta(\Lambda^2)\)) solved by
\[
y(a) = \frac{\sin(\vartheta(a)) e^{-\mathcal{H}_a^{\lambda}[\vartheta(a)]}}{\lambda \pi} \left( a f(a) e^{\mathcal{H}_a^{\lambda}[\vartheta(a)]} \cos(\vartheta(a)) \right.
\]
\[
\quad + \mathcal{H}_a^{\lambda} \left[ e^{\mathcal{H}_a^{\lambda}[\vartheta(a)]} \bullet f(\bullet) \sin(\vartheta(\bullet)) \right] + C \right) \tag{35a}
\]
\[
= \frac{\sin(\vartheta(a)) e^{\mathcal{H}_a^{\lambda}[\vartheta(a)]}}{\lambda \pi} \left( f(a) e^{-\mathcal{H}_a^{\lambda}[\vartheta(a)]} \cos(\vartheta(a)) \right.
\]
\[
\quad + \mathcal{H}_a^{\lambda} \left[ e^{-\mathcal{H}_a^{\lambda}[\vartheta(a)]} f(\bullet) \sin(\vartheta(\bullet)) \right] + \frac{C'}{\Lambda^2 - a} \right), \tag{35b}
\]
where \(\vartheta(a) = \arctan \left( \frac{\lambda \pi a}{h(a)} \right)\), \(\sin(\vartheta(a)) = \frac{|\lambda \pi|}{\sqrt{(h(a))^2 + (\lambda \pi)^2}} \geq 0\) and \(C, C'\) are arbitrary constants.

The possibility of \(C, C' \neq 0\) is due to the fact that the finite Hilbert transform has a kernel, in contrast to the infinite Hilbert transform with integration over \(\mathbb{R}\). The two formulae (35a) and (35b) are formally equivalent, but the solutions belong to different function classes and normalisation conditions may (and will) make a choice.

In principle, (35) provides the solution \(G_{ab}\) of (33), where the angle function
\[
\vartheta_b(a) := \arctan \left( \frac{\lambda \pi a}{b + \frac{1 + \lambda \pi a \mathcal{H}_a^{\lambda} [G_{\bullet 0}]}{G_{a 0}}} \right) \tag{36}
\]
plays a key rôle. This solution involves multiple Hilbert transforms which are difficult to control. A better strategy starts from the observation that the angle (36) satisfies, for \(b = 0\), again a Carleman type singular integral equation
\[
\lambda \pi \cot \vartheta_0(a) G_{a 0} - \lambda \pi \mathcal{H}_a^{\lambda} [G_{\bullet 0}] = \frac{1}{a}
\]
with solution

\[ G_{a0} = \frac{e^{-\mathcal{H}_a^\Lambda[\pi-\vartheta_0]} \sin(\vartheta_0(a))}{\lambda \pi a} \left( e^{\mathcal{H}_a^\Lambda[\pi-\vartheta_0]} \cos(\vartheta_0(a)) \right) + \mathcal{H}_a^\Lambda \left[ e^{\mathcal{H}_a^\Lambda[\pi-\vartheta_0]} \sin(\vartheta_0(\bullet)) \right] + C \) \tag{37a}

\[ = \frac{e^{\mathcal{H}_a^\Lambda[\vartheta_0]} \sin(\vartheta_0(a))}{\lambda \pi} \left( e^{-\mathcal{H}_a^\Lambda[\vartheta_0]} \cos(\vartheta_0(a)) \right) \frac{1}{a} + \mathcal{H}_a^\Lambda \left[ e^{-\mathcal{H}_a^\Lambda[\vartheta_0]} \sin(\vartheta_0(\bullet)) \right] + \frac{C'}{\lambda^2 - a}. \] \tag{37b}

Tricomi’s identities [Tri57 §4.4(28+18)], which can be arranged as

\[ e^{\pm \mathcal{H}_a^\Lambda[\vartheta_0]} \cos(\vartheta_0(a)) \mp \mathcal{H}_a^\Lambda \left[ e^{\pm \mathcal{H}_a^\Lambda[\vartheta_0]} \sin(\vartheta_0(\bullet)) \right] = 1, \]

and rational fraction expansion \( \mathcal{H}_a^\Lambda \left[ \frac{f}{\bullet} \right] = \frac{1}{a} (\mathcal{H}_a^\Lambda [f(\bullet)] - \mathcal{H}_0^\Lambda [f(\bullet)]) \) simplify (37) to

\[ G_{a0} = \frac{e^{-\mathcal{H}_a^\Lambda[\pi-\vartheta_0]} \sin(\vartheta_0(a))}{\lambda \pi a} (C - 1) \] \tag{38a}

\[ = \frac{e^{\mathcal{H}_a^\Lambda[\vartheta_0]} \sin(\vartheta_0(a))}{\lambda \pi a} \left( e^{-\mathcal{H}_a^\Lambda[\vartheta_0]} \cos(\vartheta_0(0)) + \frac{C'}{\lambda^2 - a} \right). \tag{38b} \]

Both lines are formally equivalent, but we have to guarantee the normalisation \( \lim_{a \to 0} G_{a0} = 1 \). From (36) one concludes \( \lim_{p \to 0} \vartheta_0(p) = \begin{cases} 0 & \text{for } \lambda \geq 0, \\ \pi & \text{for } \lambda < 0. \end{cases} \)

Consequently, \( e^{-\mathcal{H}_a^\Lambda[\vartheta_0]} = \exp \left( - \pi \int_0^p \frac{d\vartheta_0(p)}{\lambda^2} \right) \xrightarrow{\lambda \to 0} 0 \), which means that (38b) reduces for \( \lambda < 0 \) to (38a), with \( C' \to C - 1 \). Similarly, \( \lim_{a \to 0} e^{-\mathcal{H}_a^\Lambda[\pi-\vartheta_0]} \xrightarrow{\lambda \to 0} 0 \), so that (38a) is only consistent with \( \lambda < 0 \). The normalisation \( \lim_{a \to 0} G_{a0} = 1 \) leads with \( \lim_{a \to 0} \frac{\sin(\vartheta_0(a))}{\lambda \pi a} = 1 \) to \( 1-C = e^{-\mathcal{H}_0^\Lambda[\pi-\vartheta_0]} \) in (38a), whereas (38b) stays as it is for \( \lambda > 0 \). These results can be summarised as follows:

**Lemma 7** The angle function \( \tau_b(a) := \arctan \left( \frac{\lambda \pi a}{b + \frac{1+\lambda \pi a \mathcal{H}_a^\Lambda[\vartheta_0]}{G_{a0}}} \right) \) is for \( b = 0 \) reverted to

\[ G_{a0} = \frac{\sin(\tau_0(a))}{\lambda \pi a} \left( e^{\text{sign}(\lambda)(\mathcal{H}_a^\Lambda[\tau_0(\bullet)]-\mathcal{H}_a^\Lambda[\tau_0(\bullet)])} \right) \left( 1 - \frac{Ca}{\lambda^2 - a} \right) \] \tag{39}

where \( C \) is an arbitrary constant.

Recall that \( G_{a0} \) forms the inhomogeneity in the Carleman equation (31). We insert (39) into the Carleman solution (35) for (31) and obtain with the addition theorem \( \lambda \pi a \sin(\tau_0(a) - \tau_b(a)) = (b-d) \sin(\tau_0(a)) \sin(\tau_d(a)) \) after essentially the same steps as in the proof of (39):

18
Theorem 8 ([GW14]) The full matrix 2-point function $G_{ab}$ of self-dual $\phi^4_4$-theory on Moyal space is in the limit $\theta \to \infty$ given in terms of the boundary 2-point function $G_{a0}$ by the equation

$$G_{ab} = \frac{\sin(\tau_0(a))}{|\lambda| \pi a} e^{\text{sign}(\lambda) (\mathcal{H}_0[0\pi] - \mathcal{H}_0[\pi \pi])} \left\{ \begin{array}{ll} \frac{1}{\Lambda^2 - \alpha} & \text{for } \lambda < 0 , \\ (1 + \frac{C_a + b F(b)}{\Lambda^2 - \alpha}) & \text{for } \lambda > 0 , \end{array} \right. \quad (40)$$

where $C$ is an undetermined constant and $b F(b)$ an undetermined function of $b$ vanishing at $b = 0$.

Some remarks:

- We have proved this theorem in 2012 for $\lambda > 0$ under the assumption $C' = 0$ in [55], but knew that non-trivial solutions of the homogeneous Carleman equation parametrised by $C' \neq 0$ are possible. That no such term arises for $\lambda < 0$ (if angles are redefined $\vartheta \mapsto \tau$) is a recent result [GW14].

- We expect $C, F$ to be $\Lambda$-dependent so that $(1 + \frac{C_a + b F(b)}{\Lambda^2 - \alpha}) \xrightarrow{\Lambda \to \infty} 1 + \tilde{C} a + b \tilde{F}(b)$.

- An important observation is $G_{ab} \geq 0$, at least for $\lambda < 0$. This is a truly non-perturbative result; individual Feynman graphs show no positivity at all!

- As in [GW09], the equation for $G_{ab}$ can be solved perturbatively. Matching at $\lambda = 0$ requires $C, F$ to be flat functions of $\lambda$ (all derivatives vanish at zero). Because of $\mathcal{H}_0^\Lambda [G_{* *}] \xrightarrow{a \to \Lambda^2} -\infty$, the naïve arctan series is dangerous for $\lambda > 0$. Unless there are cancellations, we expect zero radius of convergence!

- From (40), we deduce the finite wavefunction renormalisation

$$\mathcal{Y} := -1 - \frac{dG_{ab}}{db} \bigg|_{a=b=0} = \int_0^\Lambda^2 \frac{dp}{(\lambda \pi p)^2 + \left(1 + \frac{\pi \lambda p \mathcal{H}_0^\Lambda [G_{* *}]}{G_{p0}}\right)^2} \left\{ \begin{array}{ll} 0 & \text{for } \lambda < 0 , \\ F(0) & \text{for } \lambda > 0 . \end{array} \right. \quad (41)$$

- The partition function $Z$ is undefined for $\lambda < 0$. But the Schwinger-Dyson equations for $G_{ab}$ and for higher functions, and with them log $Z$, extend to $\lambda < 0$. These extensions are unique but probably not analytic in a neighbourhood of $\lambda = 0$.

It remains to identify the boundary function $G_{a0}$. The Carleman equation (31) for $G_{ab}$ was obtained from the difference $[30]_b - [30]_{b=0}$. Consequently, $[30]_{b=0}$ gives the second relation between $G_{ab}$ and $G_{a0}$ from which both are determined. Combining them we obtain a single consistency equation for $G_{a0}$, which in terms of $T_a := |\lambda| \pi a \cot \tau_0(a)$ reads [GW12b]

$$T_a = 1 + a + \lambda \pi a \mathcal{H}_a^\Lambda [1] + \int_0^\Lambda^2 dp \frac{p \exp \left( \mathcal{H}_a^\Lambda \left[ \arctan \frac{|\lambda| \pi p}{p + T_a} \right] \right)}{\sqrt{(\lambda \pi a)^2 + (p + T_a)^2}} - \frac{p \exp \left( \mathcal{H}_0^\Lambda \left[ \arctan \frac{|\lambda| \pi p}{p + T_a} \right] \right)}{1 + p} \quad (42)$$
This equation is, unfortunately, of little use. The integrals are individually divergent for $\Lambda \to \infty$ so that we have to rely on cancellations on which we have no control.

We compensate this lack by a symmetry argument. Given the boundary function $G_{ab}$, the Carleman theory computes the full 2-point function $G_{ab}$ via (40). In particular, we get $G_{0b}$ as function of $G_{a0}$. But the 2-point function is symmetric, $G_{ab} = G_{ba}$, and the special case $b = 0$ leads to the following self-consistency equation:

$$G_{b0} = \left\{ \begin{array}{ll}
1 & \text{for } \lambda < 0 \\
1 + bF(b) & \text{for } \lambda > 0 
\end{array} \right\} \exp \left( -\lambda \int_0^b dt \int_0^{\Lambda^2} dp \left( \frac{dp}{G_{p0}(G_{0p}(G_{p0}))^2} \right) \right). \quad (43)
$$

At this point we can eventually send $\Lambda \to \infty$. Any solution of (43) is automatically smooth and (for $\lambda > 0$ but $F(b) = 0$) monotonously decreasing. Any solution of the true equation (30) (without the difference to $b = 0$) also solves the master equation (43), but not necessarily conversely. In case of a unique solution of (43) it is enough to check one candidate.

Existence of a solution of (43) is established (for $\lambda > 0$ but $F(b) = 0$) by the Schauder fixed point theorem. We consider the following subset of continuously differentiable functions on $\mathbb{R}_+$ vanishing at $\infty$:

$$K_\lambda := \left\{ f \in C_0^1(\mathbb{R}_+) : f(0) = 1, \quad 0 < f(b) \leq \frac{1}{1 + b}, \quad 0 \leq f'(b) \leq \left( \frac{1}{1 + b} + C_\lambda \right) f(b) \right\},$$

where $C_\lambda$ is defined via $2\lambda P_\lambda^2(1 + C_\lambda)e^{C_\lambda p_\lambda} = 1$ at $P_\lambda = \frac{\exp(-\frac{1}{\sqrt{1 + 4\lambda}})}{\sqrt{1 + 4\lambda}}$. Then \[GW22\]:

1. $K_\lambda$ is convex,
2. $TK_\lambda \subseteq K_\lambda$,
3. $(Tf)''(b) \leq \left( \frac{23}{4} + \frac{2}{\pi} + \frac{7 + 8\pi}{2(\lambda + 2\lambda^2)} \right) (Tf)(b)$ for any $f \in K_\lambda$,
4. $T : K_\lambda \to K_\lambda$ is continuous.

The properties 1.–3. imply that $TK_\lambda$ is relatively compact in $K_\lambda$ by a variant of the Arzelà-Ascoli theorem. Together with 4. the Schauder fixed point theorem then guarantees that (43) has a solution $G_{a0} \in K_\lambda$.

This solution provides $G_{ab}$ via (40) and all higher correlation functions via the universal algebraic recursion formulae (10), (65), (66), etc., or via the linear equations for the basic $(N_1 + \ldots + N_B)$-point functions such as (63) and (64). The
recursion formula (16) becomes after transition to continuous matrix indices
\[
G_{b_0 \ldots b_{N-1}} = \frac{(-\lambda)}{(1 + \mathcal{Y})^2} \sum_{l=1}^{N-2} G_{b_0 b_1 \ldots b_{2l-1}} G_{b_{2l} b_{2l+1} \ldots b_{N-1}} - G_{b_0 b_1 \ldots b_{2l-1}} G_{b_{2l} b_{2l+1} \ldots b_{N-1}} (b_0 - b_{2l})(b_1 - b_{N-1})
\] (44)

It involves the finite wavefunction renormalisation \(1 + \mathcal{Y} = -\frac{d G_{ab}}{d b} \bigg|_{a=b=0} \) given by (11). Of particular interest is the effective coupling constant \(\lambda_{\text{eff}} = -G_{0000}\). This limit of coinciding indices is not so easy; therefore we directly solve the integral equation for \(G_{a000}\) before using the reality condition. We find (GW12b)

\[
\lambda_{\text{eff}} = \lambda \left(1 + \frac{1}{(1 + \mathcal{Y})^2} \int_0^\infty dp \frac{\left(G_{p0} - G_{p0}\right)}{(\lambda \pi p G_{p0})^2 + (1 + \lambda \pi p \mathcal{H}_p^\infty |G_{\bullet 0}|)^2}\right).
\] (45)

The equation for the basic function \(G_{ab|cd}\) arising from (64) is solved in two steps. A first summation over \(b \in I\) in (64) yields after passage to the integral representation a Carleman equation for \(X_{a|cd} := \int_0^\infty q dq G_{aq|cd}\),

\[
X_{a|cd}\left\{1 + \lambda \int_0^\infty dq \left(\frac{G_{aq} - G_{0q}}{\lambda \pi a} - \lambda \int_0^\infty dq G_{aq} \sin \tau_q(a) \cos \left(\frac{\tau_q(a) - \tau_0(a)}{\sin \tau_0(a)}\right)\right)\right. \\
+ \mathcal{H}_a \left[X_{\bullet|\bullet} \int_0^\infty q dq \sin^2 \tau_q(\bullet) G_{aq}\right]
\]
\[
\quad = \lambda \int_0^\infty q dq (F_{aq|cd} + F_{aq|dcd}) + \frac{\lambda}{(1 + \mathcal{Y})^2} (G_{acdc} + G_{adcd}),
\]

where \(F_{ab|cd} := \frac{G_{bc|cd} - G_{bc|cd}}{G_{bc|cd} - G_{bc|cd}}\). Inserted back into (64) gives (after passage to the integral representation) a familiar Carleman equation for \(G_{ab|cd}\) with solution

\[
G_{ab|cd} = F_{ab|dcb} + F_{ab|dcd}
\]
\[
- \frac{\sin \tau_b(a)}{\lambda \pi a} \cos \tau_b(a) G_{ab} X_{a|cd} - G_{ab} \mathcal{H}_a \left[\frac{\sin^2 \tau_b(\bullet)}{\lambda \pi a} X_{\bullet|cd}\right].
\] (46)

The \(2+2\)-point function \(G_{ab|cd}\) turns out to be the most interesting part of the 4-point function in position space (see sec. 4).

3.4 Perturbation theory

The master equation (13) can, for \(F(b) \equiv 0\), be iteratively solved. To lowest order one has \(G_{a0} = \frac{1}{1+a} + \mathcal{O}(\lambda)\), from which the next order becomes

\[
G_{a0} = \frac{1}{1+a} - \lambda \log(1+a) \frac{1}{(1+a)} + \mathcal{O}(\lambda^2).
\] (47)
If we put in $G_{a0} = \frac{1}{(1+a)^{1+a}} + \mathcal{O}(\lambda^2)$ the index $a \mapsto \frac{p^2}{\mu^2}$, see (58), we get

$$\int_{\mathbb{R}^4} \frac{dp}{(2\pi\mu)^4} e^{ip(x-y)} G_{\mu,0} = \frac{2^{-\lambda}}{4\pi^2 \Gamma(1+\lambda)} \left( \frac{K_{1-\lambda}(\mu \|x-y\|)}{(\mu \|x-y\|)^{1-\lambda}} \right) \rightarrow \frac{2^{-2\lambda} \Gamma(1-\lambda)}{4\pi^2 \Gamma(1+\lambda)} \frac{1}{(\mu \|x-y\|)^{2-2\lambda}},$$

where $K_\nu(x)$ is the modified Bessel function. We thus conclude that the anomalous dimension is $\eta = -2\lambda$, i.e. negative for the stable sign $\lambda > 0$ of the coupling constant. We shall see in the next section that this result excludes a Wightman theory for $\lambda > 0$. It is worthwhile to mention that this wrong sign is a consequence of renormalisation. The divergent bare 2-point function would lead to the opposite sign. Removing the divergence at $a = 0$ overcompensates for $a > 0$ and gives $\eta = -2\lambda$. In two dimensions, $\eta$ would be non-negative for $\lambda > 0$.

From (17) we get:

- Hilbert transform: $\lambda \pi \mathcal{H}_\infty G_{0,0} = -\lambda \log(a) + \mathcal{O}(\lambda^2)$,
- angle function: $\tau_b(a) = \frac{\lambda a}{1+a+b} \left( 1 - \lambda \frac{\log(1+a) \log(1+b) - a \log a}{(1+a+b)} \right) + \mathcal{O}(\lambda^3)$,
- wavefunction renormalisation: $1 + \mathcal{Y} = -\frac{dG_{a,b}}{da}|_{a=0} = 1 + \lambda + \mathcal{O}(\lambda^2)$.

Inserted into (40) one finds

$$G_{ab} = \frac{1}{1+a+b} - \lambda \frac{(1+a) \log(1+a) + (1+b) \log(1+b)}{(1+a+b)^2} + \mathcal{O}(\lambda^2). \quad (48)$$

This result coincides with renormalised 1-loop ribbon graph computation. From the action functional (27) one obtains in the infinite volume limit to continuous matrix indices the following Feynman rules:

- $\begin{array}{c} \begin{array}{c} \begin{array}{c} a \end{array} \end{array} \end{array} \quad \begin{array}{c} \begin{array}{c} \begin{array}{c} b \end{array} \end{array} \end{array} = \frac{1}{1+(a+b)(1+\mathcal{Y})}$
- $\begin{array}{c} \begin{array}{c} \begin{array}{c} a \end{array} \end{array} \end{array} \quad \begin{array}{c} \begin{array}{c} \begin{array}{c} b \end{array} \end{array} \end{array} = -Z^2 \lambda \quad \text{(index conserved at every corner)}$
- $\begin{array}{c} \begin{array}{c} \begin{array}{c} a \end{array} \end{array} \end{array} \quad \begin{array}{c} \begin{array}{c} \begin{array}{c} b \end{array} \end{array} \end{array} = 1 + \mathcal{Y}^2 \int_0^\infty p \, dp \quad \text{for every closed face}$

To lowest order we have $G_{ab} = \frac{1}{1+(a+b)(1+\mathcal{Y}) - \Gamma_{ren}^{ab}}$, where $\Gamma_{ren}^{ab}$ is the Taylor remainder of

$$\Gamma_{ab} = \frac{p}{a} + \frac{p}{b} + \mathcal{O}(\lambda^2).$$
\begin{align}
= & \int_0^{\Lambda^2} p \, dp \, \frac{(-\lambda)}{1 + a + p} + \int_0^{\Lambda^2} p \, dp \, \frac{(-\lambda)}{1 + b + p} + \mathcal{O}(\lambda^2), \\
= & \Gamma^{\text{ren}} + \left( \int_0^{\Lambda^2} p \, dp \, \frac{(-\lambda)}{1 + p} + \int_0^{\Lambda^2} p \, dp \, \frac{(+\lambda)a}{(1 + p)^2} + (a \mapsto b) + \mathcal{O}(\lambda^2) \right), \tag{49a}
\end{align}

\Gamma_{ab}^{\text{ren}} = (-\lambda) \int_0^{\Lambda^2} p \, dp \left( \frac{1}{1 + a + p} - \frac{1}{1 + b + p} + \frac{a}{(1 + p)^2} \right) + (a \mapsto b) + \mathcal{O}(\lambda^2), \tag{49b}
in agreement with \cite{GW14}. There is no doubt that the fixed point solution for \( G_{a_0} \) and the Carleman solution for \( G_{ab} \) capture the resummation of infinitely many renormalised Feynman graphs!

From \cite{GW13} and \( \mathcal{Y} = \lambda + \mathcal{O}(\lambda^2) \) we obtain for the 4-point function

\begin{align}
G_{abcd} = & \frac{(-\lambda)}{(1 + \mathcal{Y})^2} \frac{G_{ab}G_{cd} - G_{ad}G_{cd}}{(a - c)(b - d)} =: G_{ab}G_{cd}G_{da}(-\Gamma_{abcd}), \\
\Gamma_{abcd} = & \lambda \left( 1 - \lambda \left( a - (1 + a) \log(1 + a) - c + (1 + c) \log(1 + c) \right) \right) \\
& - \lambda \left( b - (1 + b) \log(1 + b) - d + (1 + d) \log(1 + d) \right) + \mathcal{O}(\lambda^3), \tag{50}
\end{align}

which agrees with

\begin{align}
\Gamma_{abcd} = & -a \bigg( \frac{(-\lambda)}{(1 + \mathcal{Y})^2} \bigg) + \mathcal{O}(\lambda^3) \\
= & -(-\lambda) \left( 1 + 2\lambda \int_0^{\Lambda^2} p \, dp \right) + \mathcal{O}(\lambda^3) \\
= & -(-\lambda)^2 \int_0^{\Lambda^2} p \, dp \frac{1}{1 + \mathcal{Y} + \mathcal{O}(\lambda^3)} \\
& - (-\lambda)^2 \int_0^{\Lambda^2} p \, dp \frac{1}{1 + \mathcal{Y} + \mathcal{O}(\lambda^3)} \\
& - (-\lambda)^2 \int_0^{\Lambda^2} p \, dp \frac{1}{1 + \mathcal{Y} + \mathcal{O}(\lambda^3)}. \tag{51}
\end{align}

The singularities of \( Z^2 \) and of the one-loop 4-point graphs cancel exactly!

### 3.5 Computer simulations \cite{GW13}

A numerical investigation of \cite{GW13}, for \( F(b) \equiv 0 \), reveals interesting properties of the \( \phi^4 \)-theory on Moyal space. We approximate \( G_{a_0} \) as piecewise linear function on \([0, \Lambda^2]\) sampled according to a geometric progression and view \cite{GW13} as iteration \( G_{a_0}^{n+1} = (TG^n)_{a_0} \) for some initial function \( G^0 \). In this way we find numerically that \( T \) satisfies, for any \( \lambda \in \mathbb{R} \), the assumptions of the Banach fixed point theorem for Lipschitz functions on \([0, \Lambda^2]\), i.e. \( T \) is contractive and \( (G^n) \) converges to a fixed point which approximates \( G_{a_0} \). Whereas \( (G^n) \) converges for any sign of \( \lambda \).
(without discontinuity at $\lambda = 0$), the necessary consistency condition $G_{ab} = G_{ba}$ for (40) turns out to be maximally violated for $\lambda > 0$ (assuming $C = 0 = F(b)$) and satisfied (within numerical error bounds) for $\lambda \leq 0$. The observed relative asymmetry $\sup_{a,b} \left| \frac{G_{ab} - G_{ba}}{G_{ab} + G_{ba}} \right|$ of nearly 100% for $\lambda > 0$ signals that the parameters $C, F(b)$ in (40) which reflect the non-trivial solution of the homogeneous Carleman equation are definitely non-zero. Taking $C, F(b) \neq 0$ for $\lambda > 0$ into account is not feasible at the moment so that our numerical results are reliable only for $\lambda \leq 0$.

For $\lambda = 10^7$ and only 2000 sample points in $[0, \Lambda^2]$, the relative asymmetry for $\lambda \leq 0$ is of the order of 5%.

The most striking outcome of our computer simulations concerns the finite wavefunction renormalisation $(1 + \mathcal{Y})$ given by (41). Figure 1 shows both $\mathcal{Y}$ and the effective coupling constant $\lambda_{\text{eff}}$ given by (45) as functions of $\lambda$. We find clear evidence for a second-order phase transition: $\mathcal{Y}'$ is discontinuous at $\lambda_c = -0.396$, and we have in reasonable approximation a critical behaviour

$$1 + \mathcal{Y} = \begin{cases} A(\lambda - \lambda_c)\alpha & \text{for } \lambda \geq \lambda_c, \\ 0 & \text{for } \lambda < \lambda_c, \end{cases}$$

(52)

Figure 1: $\mathcal{Y}, \lambda_{\text{eff}}$ based on $G_{a0}$ for $\Lambda^2=10^7$ with 2000 sample points.
for some $A, \alpha > 0$. To be precise, we find $1 + \mathcal{Y} = 0$ only at $\lambda_0 = -0.455$, but this seems to be due to the discretisation. Of course, there cannot be a discontinuity in $\mathcal{Y}$ for finite $\Lambda$, but Figure 1 is strong support for a critical behaviour in the limit $\Lambda^2 \to \infty$. It is worthwhile to mention that nothing particular happens at the expected pole $\lambda_b = -\frac{1}{12} = 0.014$ of Borel resummation! Since $1 + \mathcal{Y} = 0$ (within numerical error bounds) in the phase $\lambda < \lambda_c$, we see from (44) that higher $N$-point functions will not exist for $\lambda < \lambda_c$. Most surprisingly, as we discuss at the end of section 4.2, a key property of the Schwinger 2-point function $S_c(x, y)$ in position space is precisely realised in $[\lambda_c, 0]$, not outside! In fact, as shown in Figure 2, one has in reasonable approximation $G_{ab} = 0$ for $0 \leq a, b \leq \Lambda^2_0$, where

$$\Lambda^2_0$$ increases with $\lambda_c - \lambda > 0$. This could leave the possibility of meaningful higher functions for matrix indices $0 \leq a_i \leq \Lambda^2_0$, but not for larger indices. Such a picture could have the interpretation of a maximal momentum cut-off of the Euclidean particles in the phase $\lambda < \lambda_c$.

### 4 Schwinger functions and reflection positivity

In the previous section we have constructed the connected matrix correlation functions $G_{[q_1^1 \cdots q_N^1 | \cdots | q_1^B \cdots q_B^B]}$ of the $(\theta \to \infty)$-limit of $\phi_4^4$-theory on Moyal space. These functions arise from the topological expansion of the free energy

$$\log \frac{Z[J]}{Z[0]} = \sum_{B=1}^{\infty} \sum_{1 \leq N_1 \leq \cdots \leq N_B} (V\mu^4)^{2-B} S_{N_1 \cdots N_B} \sum_{q_{\beta}^\beta \in \mathbb{N}^2} G_{[q_1^1 \cdots q_{N_1}^1 | \cdots | q_1^B \cdots q_B^B]} \prod_{\beta=1}^{B} \frac{1}{N_\beta^{1/2}} \left( \frac{J_{q_1^\beta} q_{2}^\beta q_{3}^\beta \cdots J_{q_{N_\beta}^\beta q_{1}^\beta}}{\mu^3} \right).$$

(53)

Since $\lim_{V\mu^4 \to \infty} G_{[q_1^1 \cdots q_{N_1}^1 | \cdots | q_1^B \cdots q_B^B]}$ is finite, the limit $\lim_{V\mu^4 \to \infty} \frac{1}{V\mu^4} \log \frac{Z[J]}{Z[0]}$ of the naturally expected free energy density removes (in addition to the removal of higher-genus contributions) all contributions from $B \geq 2$. As shown in previous sections, this planar limit is an exactly solvable and (without any doubt) non-trivial matrix model.

\[\text{Figure 2: Plots of } \log G_{a0} \text{ and } \log G_{aa} \text{ over } \log(1 + a) \text{ for } \lambda < \lambda_c.\]
4.1 Schwinger functions

We are interested here in another limit to Schwinger functions \[ \text{Sch59} \] in position space. For this end we revert the matrix representation (24) and take the infinite volume limit

\[ V \mu^4 \to \infty, \]

where we carefully have to pass to densities. Absolute position \( x \in \mathbb{R}^4 \) have no meaning, only \( \mu x \) can be used. This means that we consider (up to a factor discussed below)

\[ \langle \phi(\mu x_1) \ldots \phi(\mu x_N) \rangle \equiv \sum_{m_1, n_1, \ldots, m_N, n_N \in \mathbb{N}^2} f_{m_1 n_1}(\mu x_1) \cdots f_{m_N n_N}(\mu x_N) \langle \Phi_{m_1 n_1} \cdots \Phi_{m_N n_N} \rangle, \]

where the matrix correlation functions \( \langle \Phi_{m_1 n_1} \cdots \Phi_{m_N n_N} \rangle \) are obtained by derivatives of (53) with respect to \( J_{m_1 n_1}, \ldots, J_{m_N n_N} \). We shall see in this section that the additional index summation over \( m_i, n_i \in \mathbb{N}^2 \) gives a meaningful limit only if we redefine the volume factor in the free energy density to

\[ F[\mu] := \frac{1}{16 \pi^2 V^2 \mu^8} \log \left( \frac{\int D[\Phi] e^{-S[\Phi] + V \sum_{a \in \mathbb{N}^2} \Phi_{a a} J_a}}{\int D[\Phi] e^{-S[\Phi]}} \right) \]

\[ \mu^N S_c(\mu x_1, \ldots, \mu x_N) = \lim_{V \mu^4 \to \infty} \sum_{m_1, n_1, \ldots, m_N, n_N \in \mathbb{N}^2} f_{m_1 n_1}(\mu x_1) \cdots f_{m_N n_N}(\mu x_N) \langle \Phi_{m_1 n_1} \cdots \Phi_{m_N n_N} \rangle, \]

where \( S[\Phi] \) is given by (27) and \( f_{m n} \) by (24). By \( ( ) \) we symbolise the renormalisation of sec. 3.2.

Note that by construction the \( J \)-derivatives, and hence the Schwinger functions, are fully symmetric in \( \mu x_1, \ldots, \mu x_N \). Applying the \( J \)-derivatives the the topological expansion (53) into \( J \)-cycles produces an \( f_{m n} \)-cycle for each of the \( B \) boundary components:

\[ S_c(\mu x_1, \ldots, \mu x_N) = \lim_{V \mu^4 \to \infty} \frac{1}{64 \pi^2} \sum_{N_1 + \ldots + N_B = N} \sum_{q_{1B} \in \mathbb{N}^2} G_{q_1, \ldots, q_B} \frac{f_{q_1 q_1}(\mu x_{\sigma(1)} + \ldots + N_{\beta} + 1)}{V \mu^4 N_{\beta}} \]

\[ \times \prod_{\sigma \in S_N, \beta = 1} \frac{f_{q_2 q_2}(\mu x_{\sigma(1)} + \ldots + N_{\beta - 1}) \cdots f_{q_B q_B}(\mu x_{\sigma(1)} + \ldots + N_{\beta})}{V \mu^4 N_{\beta}}. \]
We compute the sum over the indices \( q_i^\beta \in \mathbb{N}^d \) by Laplace-Fourier transform of \( G \). For that we temporarily assume that \( \tilde{G} \) has, for every boundary component, a representation as Laplace transform in the total sum of index norms and Fourier transform in differences of index norms. This transform will be reverted in the end so that the analyticity assumption is not necessary (future analytic continuation to Minkowski space would imply representation as Laplace transform):

\[
G_{|q_1^1 \cdots q_N^i | \cdots q_N^B |} = \int_{\mathbb{R}^d} d(t^1, \ldots, t^B) \int_{\mathbb{R}^{d-B}} d(\omega^1, \ldots, \omega^i_{N_i-1}, \ldots, \omega^1_B, \ldots, \omega^B_{N_B-1})
\]

\[
\times G(t^1, \omega^1, \ldots, \omega^i_{N_i-1}| \cdots |t^B, \omega^B, \ldots, \omega^B_{N_B-1})
\]

\[
\times \prod_{\beta=1}^B \exp \left( -\frac{t^\beta}{\sqrt{V \mu^4}} \sum_{i=1}^{N_\beta} |q^i_\beta| + \frac{i}{\sqrt{V \mu^4}} \sum_{i=1}^{N_\beta-1} \omega^i (|q^i_\beta| - |q^i_{\beta+1}|) \right). \tag{55}
\]

Note that the 1-norms \(|q^i_\beta| = q^i_{\beta,1} + q^i_{\beta,2}\) imply a factorisation of the exponential, \(\exp(\ldots) = \prod_i (z^\beta_i(t^\beta, \omega^\beta))^q_{\beta,1} (\tilde{z}^\beta_i(t^\beta, \omega^\beta))^q_{\beta,2}\).

For every boundary component \( \beta = 1, \ldots, B \), we thus need to compute

\[
\sum_{q_1, \ldots, q_{N^\beta}=0}^\infty \frac{f_{q_1 q_2}(\mu \tilde{y}_1) \cdots f_{q_{N^\beta} q_{N^\beta}}(\mu \tilde{y}_{N^\beta})}{\sqrt{V \mu^4} N^{N^\beta}} \tilde{z}_1 \cdots \tilde{z}_{N^\beta}
\]

\[
= 2^{N^\beta} \sum_{q_1, \ldots, q_{N^\beta}=0}^\infty e^{-\frac{1}{2} (r_{1^\beta} + \cdots + r_{N^\beta})} \frac{L_{q_1}^{q_{1^\beta} - q_1}(r_1) \cdots L_{q_{N^\beta}}^{q_{N^\beta} - q_{N^\beta}}(r_{N^\beta})}{\sqrt{V \mu^4} N^{N^\beta}}, \tag{56}
\]

where \( r_i = \frac{\mu^2 |\tilde{y}_i|^2}{\sqrt{V \mu^4}} \) and \( \tilde{z}_j = \frac{\tilde{y}_j}{\tilde{y}_i} \exp(-\frac{t^\beta - i(\omega^\beta_1 - \omega^\beta)}{\sqrt{V \mu^4}}) \), with \( \tilde{y}_i \in \mathbb{C}, \tilde{y}_0 \equiv \tilde{y}_{N^\beta} \) and \( \omega_0 = \omega_{N^\beta} \equiv 0 \). One has

**Lemma 11 ([GW13b])** For \(|\tilde{z}_j| < 1 \), a cyclic product of Laguerre polynomials (i.e. \( N^\beta + j \equiv j \)) is summed to

\[
\sum_{q_1, \ldots, q_{N^\beta}=0}^\infty \prod_{j=1}^{N^\beta} (\tilde{z}_j)^{q_j} L_{q_j}^{q_{j+1} - q_j}(r_j) = \frac{\exp \left( \sum_{j,k=1}^{N^\beta} r_j (-\tilde{z}_j^{k+j}) \cdots (-\tilde{z}_{N^\beta}^{N+j}) \right)}{1 - (-\tilde{z}_1) \cdots (-\tilde{z}_{N^\beta})}. \tag{57}
\]

The denominators in (57) become

\[
1 - (-\tilde{z}_1) \cdots (-\tilde{z}_{N^\beta}) = 1 - (-1)^{N^\beta} \exp \left( \frac{-N^\beta t}{\sqrt{V \mu^4}} \right) \overset{V \mu^4 \to \infty}{\longrightarrow} \begin{cases} 
\frac{N^\beta t}{\sqrt{V \mu^4}} & \text{for } N^\beta \text{ even,} \\
\frac{N^\beta t}{2 \sqrt{V \mu^4}} & \text{for } N^\beta \text{ odd.}
\end{cases}
\]

Together with the prefactor \( \frac{1}{\sqrt{V \mu^4}} \), the sum (55) converges for \( V \mu^4 \to \infty \) to zero if \( N^\beta \) is odd, whereas if \( N^\beta \) is even the limit is non-zero and finite, depending
only on $t$ but no longer on $\omega_j$. Recombining the two $N^2$-components we produce factors
\[
\exp \left( -\frac{|\mu x|^2}{2Nt} \right) = \int_{\mathbb{R}^d} \frac{dp}{4\pi^2\mu^4} e^{-i\langle \frac{2}{\mu} \mu x \rangle} \exp \left( -\frac{N't\|p\|^2}{2\mu^2} \right)
\] for every even $N'$.

Altogether we arrive at
\[
\lim_{V\mu^4 \to \infty} \sum_{q_1,\ldots,q_{N'} \in \mathbb{N}^d} f_{q_1\ldots q_{N'}}(\mu x_1) \cdots f_{q_1,\ldots,q_{N'}}(\mu x_{N'}) \sum_{q_{N',1}+q_{N',2} \cdots q_{N',1}q_{N',2}}
\]
\[
= \frac{4^{N'}}{N'} \int_{\mathbb{R}^d} \frac{dp}{4\pi^2\mu^4} e^{-i\langle \frac{2}{\mu} \mu x_1 - x_2 + \cdots + x_{N'-1} - x_{N'} \rangle} \exp \left( -\frac{N't\|p\|^2}{2\mu^2} \right)
\] for $N'$ even,
\[
= 0
\] for $N'$ odd.

Integration of $G(t^1, \omega_1^1, \ldots, \omega_{N_1-1}^1; \ldots; t^B, \omega_1^B, \ldots, \omega_{N_B-1}^B)$ against $\exp(-\frac{N't\|p\|^2}{2\mu^2})$ in (55) returns to the original function $G|_{q_1^1\ldots q_{N_1}^1;\ldots; q_B^1\ldots q_{N_B}^1}$, but with

1. for each $\beta$, all $|q_{\beta}^j|$ coincide (no $\omega$-dependence),
2. $\frac{\sum_{j=1}^{N_{\beta}} |q_{\beta}^j|}{\sqrt{V\mu^4}} = \frac{N_{\beta}}{2\mu^4}$, hence $\frac{|q_{\beta}^j|}{\sqrt{V\mu^4}} \to (1+\gamma)q = \frac{|p|^2}{2\mu^2}$ in the limit to the integral representation.

We have thus proved (GW13b):

**Theorem 12** The connected $N$-point Schwinger functions of the $\phi_4^2$-model on extreme Moyal space $\theta \to \infty$ are given by

\[
S_c(\mu x_1, \ldots, \mu x_N)
\]
\[
= \frac{1}{64\pi^2} \sum_{N_1,\ldots,N_B=N} \sum_{\sigma \in S_N} \left( \prod_{\beta=1}^{B} \frac{4^{N_\beta}}{N_\beta} \int_{\mathbb{R}^d} \frac{dp_{\beta}}{4\pi^2\mu^4} e^{i\langle \sum_{i=1}^{N_\beta} (-1)^{i-1} \mu x_{\sigma(N_1+\ldots+N_{\beta-1}+i)} \rangle} \right)
\]
\[
\times G_{\begin{array}{c} p_1^1 \|p_1^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{\|p_1^1\|^2}{2\mu^2(1+\gamma)} \\
\|p_1^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \end{array}}^{N_1} \ldots
\]
\[
\begin{array}{c} p_B^1 \|p_B^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \\
\|p_B^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \end{array}^{N_B}.
\]

(58)

Some comments:

- Only a restricted sector of the underlying matrix model contributes to position space: All strands of the same boundary component carry the same matrix index.
- Schwinger functions are symmetric and invariant under the full Euclidean group. This comes truly surprising since $\theta \neq 0$ breaks both translation invariance and manifest rotation invariance. The limit $\theta \to \infty$ was expected to make this symmetry violation even worse!
- The most interesting sector is the case where every boundary component has $N_\beta = 2$ indices. It is described by the $(2+\ldots+2)$-point functions $G_{\begin{array}{c} p_1^1 \|p_1^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{\|p_1^1\|^2}{2\mu^2(1+\gamma)} \\
\|p_1^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \end{array}}^{N_1} \ldots
\]
\[
\begin{array}{c} p_B^1 \|p_B^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \\
\|p_B^1\|^2 \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \ldots \frac{2\mu^2(1+\gamma)}{2\mu^2(1+\gamma)} \end{array}^{N_B}.
\]
This sector describes the propagation and interaction of $B$ particles without any momentum exchange. This is acceptable for a 2D-model. In four dimensions, absence of momentum transfer is a sign of triviality.

However, typical triviality proofs rely on clustering, analyticity in Mandelstam representation or absence of bound states. All this needs verification.

It is already clear that clustering is maximally violated. Looking for instance at the $(2+2)$-sector, we have

$$\lim_{\mu a \to \infty} S_{c}^{2+2}(\mu x_1, \mu x_2, \mu(x_3 + a), \mu(x_4 + a))$$

$$= \int \frac{dp\,dq}{4\pi^6\mu^4} G_{\frac{||p||^2}{2\mu^2(1+\gamma)} \frac{||q||^2}{2\mu^2(1+\gamma)} \frac{||q||^2}{2\mu^2(1+\gamma)}} e^{i(p\cdot(x_1-x_2)+i(q\cdot(x_3-x_4))}$$

independent of the distance between $\{x_1, x_2\}$ on one hand and $\{x_3, x_4\}$ on the other hand. Absence of clustering means that the vacuum state (of a hypothetical continuation to a Wightman theory) is not a pure state. Non-pure states can be decomposed into pure states which describe different topological sectors.

Let us give an intuitive explanation why the limit $\theta \to \infty$ of extreme non-commutativity is so close to an ordinary field theory expected for $\theta \to 0$. The interaction term in momentum space

$$\frac{\lambda}{4} \int_{(\mathbb{R}^4)^4} \left( \prod_{i=1}^{4} \frac{dp_i}{(2\pi)^4} \right) \delta(p_1 + \cdots + p_4) \exp \left( i \sum_{i<j} \langle p_i, \Theta p_j \rangle \right) \prod_{i=1}^{4} \hat{\phi}(p_i)$$

leads to the Feynman rule $\lambda \exp \left( i \sum_{i<j} \langle p_i, \Theta p_j \rangle \right)$, plus momentum conservation. For $\theta \to \infty$, this converges to zero almost everywhere by the Riemann-Lebesgue lemma, unless $p_i, p_j$ are linearly dependent. This case of linearly dependent momenta might be protected for topological reasons, and these are precisely the boundary components $B > 1$ which guarantee full Lebesgue measure!

### 4.2 Reflection positivity

Under conditions identified by Osterwalder-Schrader [OS73, OS75], Schwinger functions [Sch59] of a Euclidean quantum field theory permit an analytical continuation to Wightman functions [Wig56, SW64] of a true relativistic quantum field theory. In simplified terms, the reconstruction theorem of Osterwalder-Schrader for a field theory on $\mathbb{R}^d$ says:

**Theorem 13** ([OS73, OS75]) Assume the Schwinger functions $S(x_1, \ldots, x_N)$ satisfy

0. growth conditions,
1. Euclidean covariance,
2. reflection positivity: for each tuple \((f_0, \ldots, f_K)\) of test functions \(f_N \in S(\mathbb{R}^N)\),
\[
\sum_{M,N=0}^K \int dx \, dy \, S(x_1, \ldots, x_N, y_1, \ldots, y_M) f_N(x_1^r, \ldots, x_N^r) f_M(y_1, \ldots, y_M) \geq 0 ,
\]
where \((x^0, x^1, \ldots, x^{d-1})^r := (-x^0, x^1, \ldots, x^{d-1})\).

3. permutation symmetry.
Then the \(S(\xi_1, \ldots, \xi_{N-1})|_{\xi^0>0}\), with \(\xi_i = x_i-x_{i+1}\), are Laplace-Fourier transforms of
Wightman functions in a relativistic quantum field theory. If in addition the \(S(x_1, \ldots, x_N)\) satisfy
4. clustering
then the Wightman functions satisfy clustering, too.

Representation as Laplace transform in \(\xi^0\) requires analyticity in \(\text{Re}(\xi^0) > 0\).
For the Schwinger 2-point function (58), such analyticity in \(\xi^0\) is a corollary of
analyticity of the function \(a \mapsto G_{aa} \in \mathbb{C} \setminus [-\infty, 0]\). We will show that analyticity
and reflection positivity boil down to Stieltjes functions, i.e. functions \(f : \mathbb{R}_+ \to \mathbb{R}\)
which have a representation as a Stieltjes transform (see [Wid38])
\[
f(x) = c + \int_0^\infty \frac{d(\rho(t))}{x+t} , \quad c = f(\infty) \geq 0 , \quad (60)
\]
where \(\rho\) is non-negative and non-decreasing. We prove:

**Proposition 14** The Schwinger function \(S_c(\mu \xi) = \int_{\mathbb{R}^4} \frac{dp}{(2\pi \mu)^4} e^{ip\xi} G_{\mu^2(1+\gamma)} \)
identified in (58) is the analytic continuation of a Wightman 2-point function if
and only if \(a \mapsto G_{aa}\) is Stieltjes.

**Proof.** This is verified by explicit calculation. If \(a \mapsto G_{aa}\) is Stieltjes, we have in terms of \(\omega_p(t) := \sqrt{p^2 + 2\mu^2(1+\gamma)}\)
\[
S_c(\mu \xi)|_{\xi^0>0} = \int_{\mathbb{R}^3} \frac{dp}{(2\pi \mu)^3} \int_0^\infty \frac{dp^0}{2\pi \mu} e^{ip\xi^0 + ip\xi} \int_0^\infty \frac{dp(t)}{t + \frac{(p^0)^2 + p^2}{2\mu^2(1+\gamma)}}
\]
\[
= 2\mu(1+\gamma) \int_{\mathbb{R}^3} \frac{dp}{(2\pi \mu)^3} \int_0^\infty \frac{dp^0}{2\omega_p(t)} \int_0^\infty \frac{dp(t)}{2\pi i} \left( \frac{e^{ip\xi^0}}{p^0 - i\omega_p(t)} - \frac{e^{ip\xi^0}}{p^0 + i\omega_p(t)} \right)
\]
\[
= 2\mu(1+\gamma) \int_0^\infty \frac{dp(t)}{2\omega_p(t)} \int_{\mathbb{R}^3} \frac{dp}{(2\pi \mu)^3} \left( e^{-p\xi^0 - i\vec{p}\vec{x}} \right)
\]
\[
= \int_0^\infty \frac{2(1+\gamma)}{\mu^4} \frac{dp(t)}{dp(t)} \frac{d\hat{W}_i(q)e^{-q_i^0 + i\vec{q}\vec{x}}}{d\vec{q}} , \quad (61a)
\]
\[
\hat{W}_i(q) := \frac{\theta(q^0)}{(2\pi)^3} \delta \left( \frac{(q^0)^2 - q^2 - 2\mu^2(1+\gamma)t}{\mu^2} \right) . \quad (61b)
\]
The step from the second to third line is the residue theorem. We observe that \( W_t(q) \) is precisely the Källén-Lehmann spectral representation \([\text{Kal}52, \text{Leh}54]\) of a Wightman 2-point function.

Remarkably, the Stieltjes property can be tested by purely real conditions:

**Theorem 15 (Widder [Wid38])** A function \( f : \mathbb{R}_+ \to \mathbb{R} \) is Stieltjes iff it is smooth, non-negative and satisfies

\[
L_{k,t}[f(\bullet)] := \frac{(-t)^{k-1}}{c_k} \frac{d^{2k-1}}{dt^{2k-1}}(t^k f(t)), \quad c_1 = 1, \; c_{k>1} = k!(k-2)!
\]

In that case, the measure is recovered by \( \rho'(t) = \lim_{k \to \infty} L_{k,t}[f(\bullet)] \) (weakly and almost everywhere).

The perturbatively established anomalous dimension \( \eta = -2\lambda \) implies that \( a \mapsto G_{aa} \) cannot be Stieltjes for \( \lambda > 0 \). The restriction to negative coupling constant is reminiscent of the planar wrong-sign \( \lambda \phi^4_4 \)-model \([\text{tHo}82, \text{Riv}83]\). Recall that our matrix model also reduces to the planar sector, but as result of the infinite volume limit and not by hand. We nonetheless keep a non-trivial topology in form of \( B \geq 1 \) boundary components. Moreover, we have an exact solution for \( S(x_1, \ldots, x_N) \), not only an existence proof.

![Figure 3: Widder's criteria](image)

- based on interpolation of discrete data, noisy for \( k \geq 4 \)
- Stieltjes property clearly violated for \( \lambda < \lambda_c \)

Whether or not \( a \mapsto G_{aa} \) is a Stieltjes function for \( \lambda < 0 \) is a highly interesting question. A first idea can be obtained by computer simulations, see sec. 3.5.
show in Figure 3 interpolation results for $\lambda$ near the critical coupling constant. We find clear evidence that $a \mapsto G_{aa}$ is not a Stieltjes function for $\lambda < \lambda_c$, where $\lambda_c \approx -0.396$ locates the discontinuity of $Y'(\lambda)$. For $\lambda \in [\lambda_c, 0]$ the results are not conclusive (as $k$ is too small). Since $G_{aa}$ and $G_{a0}$ show a very similar behaviour (see e.g. Fig. 2), the functions $L_{k,t}[G_{a*}]$ (which are easy to compute) give some indication about $L_{k,t}[G_{aa}]$ (which we are interested in). From (43) one can prove the following identity [GW14]:

$$\frac{(\log G_{a0})^{(\ell)}}{(\ell - 1)!} = \frac{(-1)^{\ell}}{(1 + \alpha)^{\ell}} + (-1)^{\ell} \text{sign}(\lambda) \mathcal{H}_0^\Lambda \left[ \sin (\ell \tau_a(\bullet)) \left( \frac{|\sin \tau_a(\bullet)|}{|\lambda|} \right)^{\ell} \right].$$ (62)

The resulting integrated ‘mass densities’ $\tilde{\rho}_k(m^2) = \int_0^{m^2} dt \ L_{k,t}[G_{a*}]$ are shown in Figure 4. We find clear evidence for a mass gap, $\lim_{k \to \infty} \tilde{\rho}_k(\mu^2) = 0$ for

$$\begin{align*}
\lambda &= -0.032 \\
\lambda &= -0.366
\end{align*}$$

![Figure 4: $\tilde{\rho}_k(m^2) = \int_0^{m^2} dt \ L_{k,t}[G_{a*}]$ as approximation for the mass density of $a \mapsto G_{aa}$. In each row, the left picture is zoomed into small $\mu$, showing evidence for a mass gap. The right pictures show the global behaviour, close to a step function expected for $\lambda = 0$ close to critical behaviour: $\tilde{\rho}(\mu^2)=\begin{cases} 0 & \mu \leq m \\ (\mu^2-m^2)^{-\alpha} & \mu \geq m \end{cases}$]

$0 \leq \mu^2 \leq m^2$. For $\lambda > 0$ the integrated mass density approaches (as expected) a step function, whereas for $\lambda \lesssim \lambda_c$ we notice a power-law behaviour typical for critical phenomena. In particular, for $\lambda_c < \lambda < 0$ there is no further gap in the support of $\tilde{\rho}'$, which signals scattering right away from $m^2$ (not only from the two-particle threshold on). We interpret this as scattering of a massive particle with an infrared cloud. This scattering would be a remnant of the underlying non-trivial matrix model before the projection to diagonal matrices.
4.3 Summary

We have shown that the $\phi^4_2$-model on noncommutative Moyal space, considered in the limit $\theta \to \infty$ of extreme noncommutativity, is an exactly solvable and non-trivial matrix model. Euclidean symmetry is violated in the beginning, but we identified a limit which projects to diagonal matrices where Euclidean symmetry is restored. One would not expect that such a brutal projection can respect any quantum field theory axioms. Surprisingly, the first consistency checks, positivity of the lowest Widder criteria $L_{k,t}[G_{\bullet\bullet}]$, are passed for the only interesting interval $[\lambda_c, 0]$ of the coupling constant!

If these miracles continue and all Osterwalder-Schrader axioms (except for clustering) hold, we would get a relativistic quantum field theory in four dimensions. This theory is somewhat strange as ‘particles’ keep their momenta in interaction processes. Nevertheless, the theory is not completely trivial. We find scattering remnants from the noncommutative geometrical (i.e. matricial) substructure. Only the external matrix indices are put ‘on-shell’, internally all degrees of freedom contribute.

We have seen that clustering is maximally violated. The interaction is insensitive to positions in different boundary components. In particular, ‘particles’ are never asymptotically free.

A Schwinger-Dyson equations for $B = 2$

We find for the $(1+1)$- and $(2+2)$-point functions

$$G_{[a|c]} = -\frac{\lambda}{E_a+E_a} \left( \frac{1}{V} \sum_{p \in I} \left( G_{[ap]} G_{[a|c]} - \frac{G_{[pc]}-G_{[a|c]}}{E_p-E_a} \right) - \frac{G_{[cc]}-G_{[ac]}}{E_c-E_a} \right)$$  \hspace{1cm}(63a)

$$- \frac{\lambda}{V^2(E_a+E_a)} \left( 3G_{[a|a]} G_{[a|c]} + G_{[a|aca]} + G_{[aca]} + \frac{1}{V} \sum_{n \in I} G_{[a|an]} \right)$$  \hspace{1cm}(63b)

$$- \frac{\lambda}{V^4(E_a+E_a)} G_{[a|a|a|c]}$$  \hspace{1cm}(63c)

$$G_{[ab|cd]} = -\frac{\lambda}{E_a+E_b} \left( \frac{1}{V} \sum_{p \in I} \left( (G_{[ap]} G_{[ab|cd]} + G_{[ab|ap]} G_{[cd]}) - \frac{G_{[pb|cd]}-G_{[ab|cd]}}{E_p-E_a} \right) \right)$$  \hspace{1cm}(64a)

$$+ G_{[ab]} \left( G_{[car]} + G_{[dcar]} \right) - \frac{G_{[dce]}-G_{[cd]}}{E_c-E_a} - \frac{G_{[dbde]}-G_{[db]}-G_{[dbac]}}{E_d-E_a}$$  \hspace{1cm}(64b)

$$- \frac{\lambda}{V^2(E_a+E_b)} \left( G_{[a|a]} G_{[ab|cd]} + G_{ab} G_{[a|cd]} + \frac{1}{V} \sum_{n \in I} G_{[an|ab|cd]} \right)$$  \hspace{1cm}(64b)

$$+ G_{[cd|aba]} + G_{[cd|bab]} + G_{[ac|ca]} + G_{[ac|cd]} + G_{[ac|cd]} + G_{[ac|cd]} - \frac{G_{[b|a|cd]}-G_{[b|b|cd]}}{E_b-E_a}$$  \hspace{1cm}(64b)

$$- \frac{\lambda}{V^4(E_a+E_b)} G_{[a|a|ab|cd]}$$  \hspace{1cm}(64c)
These are basic functions which are not simplified by reality. As before, (63a) and (63b) preserve the genus, whereas \( g \mapsto g+1 \) in (63a)+(64a) and \( g \mapsto g+2 \) in (63c)+(64c). The higher \((N_1+N_2)\)-point functions with one \( N_i \geq 3 \) simplify by reality to universal recursion formulae. For \( N_i \) odd we have

\[
G_{[b_0\ldots b_{2l-1}]_{c_1\ldots c_{N_{-2l-1}}}} - \lambda \sum_{k=1}^{N-2l-1} G_{[c_1\ldots c_{k-1}b_0b_1\ldots b_{2l}c_{k+1}\ldots c_{N_{-2l-1}}]} - G_{[c_1\ldots c_{k-1}b_0b_1\ldots b_{2l}b_0c_{k+1}\ldots c_{N_{-2l-1}}]} (E_{b_0} - E_{b_{2l}})(E_{b_0} - E_{c_k})
\]

\[
- \lambda \sum_{j=1}^{l} G_{[b_0b_1\ldots b_{2j-2}]_{c_1\ldots c_{N_{-2l-1}}}} G_{[b_{2j-1}b_{2j}b_{2l}]} - G_{[b_{2j-1}b_{2j}b_{2l}]} G_{[b_0b_2\ldots b_{2l}]} (E_{b_0} - E_{b_{2j}})(E_{b_0} - E_{b_{2l}})
\]

\[
- \lambda \sum_{j=1}^{l} G_{[b_0b_1\ldots b_{2j-1}] G_{[b_{2j+1}b_{2j}b_{2l}]} - G_{[b_{2j+1}b_{2j}]} G_{[b_0b_{2j+1}\ldots b_{2l}]} (E_{b_0} - E_{b_{2j+1}})(E_{b_0} - E_{b_{2j}})
\]

\[
- \frac{\lambda}{V^2} \sum_{k=1}^{2l} G_{[b_0b_1\ldots b_{k-1}b_{2k+1}\ldots b_{2l}c_1\ldots c_{N_{-2l-1}}]} - G_{[b_0b_1\ldots b_{k-1}b_{2k+1}\ldots b_{2l}]} (E_{b_0} - E_{b_{2l}})(E_{b_0} - E_{b_{2k+1}}) \quad (65)
\]

The last line increases the genus and is absent in \( G^{(0)}_{[b_0b_1\ldots b_{2l}c_1\ldots c_{N_{-2l-1}}]} \). For \( N_i \) even one finds

\[
G_{[ab_1\ldots b_{2l-1}]_{c_1\ldots c_{N_{-2l}}}} - \lambda \sum_{j=1}^{l-1} G_{[b_1\ldots b_{2j-1}a][c_1\ldots c_{N_{-2j}}]} G_{[b_{2j+1}b_{2j+1}\ldots b_{2l-1}]} - G_{[b_{2j}][c_1\ldots c_{N_{-2j}}]} G_{[a_{2j+1}b_{2j+1}\ldots b_{2l-1}]} (E_{b_0} - E_{b_{2j+1}})(E_{a} - E_{b_{2j}})
\]

\[
- \lambda \sum_{j=1}^{l-1} G_{[b_1\ldots b_{2j-1}a]} G_{[b_{2j+1}b_{2j+1}\ldots b_{2l-1}[c_1\ldots c_{N_{-2j}}]} - G_{[b_{2j-1}b_{2j}] G_{[a_{2j+1}b_{2j+1}\ldots b_{2l-1}][c_1\ldots c_{N_{-2j}}]} (E_{b_0} - E_{b_{2j+1}})(E_{a} - E_{b_{2j}})
\]

\[
- \lambda \sum_{k=1}^{N-2l} G_{[c_1\ldots c_{k-1}ab_1\ldots b_{2l-1}c_{k+1}\ldots c_{N_{-2l}}]} - G_{[c_1\ldots c_{k-1}b_0b_1\ldots b_{2l-1}c_{k+1}\ldots c_{N_{-2l}}]} (E_{b_0} - E_{b_{2l-1}})(E_{a} - E_{b_{k}})
\]

\[
- \lambda \frac{\lambda}{V^2} \sum_{k=1}^{2l} G_{[b_0b_1\ldots b_{k-1}b_{2k+1}\ldots b_{2l}[c_1\ldots c_{N_{-2k}}]} - G_{[b_0b_1\ldots b_{k-1}b_{2k+1}\ldots b_{2l}][c_1\ldots c_{N_{-2k}}]} (E_{b_0} - E_{b_{2l}})(E_{b_0} - E_{b_{2k+1}}) \quad (66)
\]

Again, the last line increases the genus and is absent in \( G^{(0)}_{[b_0b_1\ldots b_{2l-1}c_1\ldots c_{N-2l}]} \).

**Acknowledgements**

RW would like to cordially thank Keiichi R. Ito for the invitation to the RIMS symposium “Applications of RG Methods in Mathematical Sciences” and for the immense help and hospitality during this visit to Kyoto in September 2013.

In addition to the presentation at the RIMS symposium, this contribution is based on a lecture series which both of us gave in November 2013 in Göttingen.
We would like to thank Dorothea Bahns for this arrangement and for hospitality during this series.

References

[Aiz81] M. Aizenman, “Proof of the triviality of $\phi^4_4$ field theory and some mean field features of Ising models for $d > 4$,” Phys. Rev. Lett. 47 (1981) 1–4.

[ADJ97] J. Ambjørn, B. Durhuus and T. Jonsson, Quantum geometry. A statistical field theory approach, Cambridge Univ. Press (1997).

[BDFP02] D. Bahns, S. Doplicher, K. Fredenhagen and G. Piacitelli, “On the unitarity problem in space-time noncommutative theories,” Phys. Lett. B 533 (2002) 178–181 [hep-th/0201222].

[Bah10] D. Bahns, “The ultraviolet infrared mixing problem on the noncommutative Moyal space,” arXiv:1012.3707 [hep-th].

[Bon12] V. Bonzom, “Revisiting random tensor models at large $N$ via the Schwinger-Dyson equations,” JHEP 1303 (2013) 160 [arXiv:1208.6216 [hep-th]].

[Car22] T. Carleman, “Sur la résolution de certaines équations intégrales,” Arkiv for Mat., Astron. och Fysik 16 (1922), 19pp.

[CC96] A. H. Chamseddine and A. Connes, “The spectral action principle,” Commun. Math. Phys. 186 (1997) 731–750 [hep-th/9606001].

[CR00] I. Chepelev and R. Roiban, “Convergence theorem for noncommutative Feynman graphs and renormalization,” JHEP 0103 (2001) 001 [hep-th/0008090].

[Con94] A. Connes, Noncommutative geometry, Academic Press, San Diego (1994).

[Con96] A. Connes, “Gravity coupled with matter and foundation of noncommutative geometry,” Commun. Math. Phys. 186 (1996) 155–176 [hep-th/9603053].

[DGZ95] P. Di Francesco, P. H. Ginsparg and J. Zinn-Justin, “2D gravity and random matrices,” Phys. Rept. 254 (1995) 1–133 [hep-th/9306153].

[DR07] M. Disertori and V. Rivasseau, “Two and three loops beta function of non commutative $\phi^4_d$ theory,” Eur. Phys. J. C 50 (2007) 661–671 [hep-th/0610224].

[DGMR07] M. Disertori, R. Gurau, J. Magnen and V. Rivasseau, “Vanishing of beta function of non commutative $\phi^4_1$ theory to all orders,” Phys. Lett. B 649 (2007) 95–102 [hep-th/0612251].

[DFR95] S. Doplicher, K. Fredenhagen and J. E. Roberts, “The Quantum structure of space-time at the Planck scale and quantum fields,” Commun. Math. Phys. 172 (1995) 187–220 [hep-th/9303037].

[Filk6] T. Filk, “Divergencies in a field theory on quantum space,” Phys. Lett. B 376 (1996) 53–58.

[Frö82] J. Fröhlich, “On the triviality of $\lambda \phi^4_4$ theories and the approach to the critical point in $d \geq 4$ dimensions,” Nucl. Phys. B 200 (1982) 281–296.
[GGISV03] V. Gayral, J. M. Gracia-Bondía, B. Iochum, T. Schücker and J. C. Várilly, “Moyal planes are spectral triples,” Commun. Math. Phys. 246 (2004) 569–623 [hep-th/0307241].

[GW13a] V. Gayral and R. Wulkenhaar, “Spectral geometry of the Moyal plane with harmonic propagation,” J. Noncommut. Geom. 7 (2013) 939–979 [arXiv:1108.2184 [math.OA]].

[GJ87] J. Glimm and A. M. Jaffe, Quantum physics. A functional integral point of view, Springer-Verlag, New York (1987).

[GV88] J. M. Gracia-Bondía and J. C. Várilly, “Algebras of distributions suitable for phase space quantum mechanics. I.,” J. Math. Phys. 29 (1988) 869–879.

[GW04] H. Grosse and R. Wulkenhaar, “The β-function in duality-covariant noncommutative φ⁴-theory,” Eur. Phys. J. C 35 (2004) 277–282 [hep-th/0402093].

[GW05a] H. Grosse and R. Wulkenhaar, “Power-counting theorem for non-local matrix models and renormalisation,” Commun. Math. Phys. 254 (2005) 91–127 [hep-th/0305066].

[GW05b] H. Grosse and R. Wulkenhaar, “Renormalisation of φ⁴-theory on noncommutative ℝ⁴ in the matrix base,” Commun. Math. Phys. 256 (2005) 305–374 [hep-th/0401128].

[GW05c] H. Grosse and R. Wulkenhaar, “Renormalization of φ⁴-theory on noncommutative ℝ⁴ to all orders,” Lett. Math. Phys. 71 (2005) 13–26 [hep-th/0403232].

[GW09] H. Grosse and R. Wulkenhaar, “Progress in solving a noncommutative quantum field theory in four dimensions,” arXiv:0909.1389 [hep-th].

[GW12a] H. Grosse and R. Wulkenhaar, “8D-spectral triple on 4D-Moyal space and the vacuum of noncommutative gauge theory,” J. Geom. Phys. 62 (2012) 1583–1599 [arXiv:0709.0095 [hep-th]].

[GW12b] H. Grosse and R. Wulkenhaar, “Self-dual noncommutative φ⁴-theory in four dimensions is a non-perturbatively solvable and non-trivial quantum field theory,” arXiv:1205.0465 [math-ph], to appear in Commun. Math. Phys.

[GW13b] H. Grosse and R. Wulkenhaar, “Solvable limits of a 4D noncommutative QFT,” arXiv:1306.2816 [math-ph].

[GW14] H. Grosse and R. Wulkenhaar, “Phase transitions in exactly solvable 4D noncommutative QFT,” in preparation.

[Gur10] R. Gurau, “The 1/N expansion of colored tensor models,” Annales Henri Poincare 12 (2011) 829–847 [arXiv:1011.2726 [gr-qc]].

[GP12] R. Gurau and J. P. Ryan, “Colored tensor models - a review,” SIGMA 8 (2012) 020 [arXiv:1109.4812 [hep-th]].

[JW00] A. M. Jaffe and E. Witten, “Quantum Yang-Mills theory,” in: The millennium prize problems, pp. 129–152, eds. J. Carlson et al, Amer. Math. Soc., Providence (2006).
[Käl52] G. Källén, “On the definition of the renormalization constants in quantum electrodynamics,” Helv. Phys. Acta 25 (1952) 417–434.

[Kon82] M. Kontsevich, “Intersection theory on the moduli space of curves and the matrix Airy function,” Commun. Math. Phys. 147 (1992) 1–23.

[LAK54a] L. D. Landau, A. A. Abrikosov and I. M. Khalatnikov, “On the removal of infinities in quantum electrodynamics,” (in russ.) Dokl. Akad. Nauk SSSR 95 (1954) 497–500.

[LAK54b] L. D. Landau, A. A. Abrikosov and I. M. Khalatnikov, “Asymptotic expression of the electron Green function in quantum electrodynamics,” (in russ.) Dokl. Akad. Nauk SSSR 95 (1954) 773–776.

[LAK54c] L. D. Landau, A. A. Abrikosov and I. M. Khalatnikov, “An asymptotic expression for the photon Green function in quantum electrodynamics,” (in russ.) Dokl. Akad. Nauk SSSR 95 (1954) 1117–1120.

[LS02] E. Langmann and R. J. Szabo, “Duality in scalar field theory on non-commutative phase spaces,” Phys. Lett. B 533 (2002) 168–177 [hep-th/0202039].

[Leh54] H. Lehmann, “Über Eigenschaften von Ausbreitungsfunktionen und Renormierungskonstanten quantisierter Felder,” Nuovo Cim. 11 (1954) 342–357.

[MVS00] S. Minwalla, M. Van Raamsdonk and N. Seiberg, “Noncommutative perturbative dynamics,” JHEP 0002 (2000) 020 [hep-th/9912072].

[OS73] K. Osterwalder and R. Schrader, “Axioms for Euclidean Green’s functions,” Commun. Math. Phys. 31 (1973) 83–112.

[OS75] K. Osterwalder and R. Schrader, “Axioms for Euclidean Green’s functions II,” Commun. Math. Phys. 42 (1975) 281–305.

[Rie93] M. A. Rieffel, “Deformation quantization for actions of $\mathbb{R}^d$, Memoirs AMS 506 (1993) 1–96.

[Riv83] V. Rivasseau, “Construction and Borel summability of planar four-dimensional Euclidean field theory,” Commun. Math. Phys. 95 (1984) 445–486.

[Riv91] V. Rivasseau, From perturbative to constructive renormalization, Princeton Univ. Press (1991).

[Riv07a] V. Rivasseau, “Non-commutative renormalization” In: Quantum spaces (Séminaire Poitcaré X), pp 19–109, eds. B. Duplantier and V. Rivasseau, Birkhäuser Verlag Basel (2007) [arXiv:0705.0705 [hep-th]].

[Riv07b] V. Rivasseau, “Constructive matrix theory,” JHEP 0709 (2007) 008 [arXiv:0706.1224 [hep-th]].

[Riv13] V. Rivasseau, “The Tensor Track, III,” arXiv:1311.1461 [hep-th].

[Sam14] D. O. Samary, “Closed equations of the two-point functions for tensorial group field theory,” arXiv:1401.2096 [hep-th].

[Sch59] J. Schwinger, “Euclidean quantum electrodynamics,” Phys. Rev. 115 (1959) 721–731.
| Reference | Author(s) and Title |
|-----------|--------------------|
| [SW64]    | R. F. Streater and A. S. Wightman, *PCT, spin and statistics, and all that*, Benjamin, New York (1964). |
| [tHo74]   | G. ’t Hooft, “A planar diagram theory for strong interactions,” Nucl. Phys. B **72** (1974) 461–473. |
| [tHo82]   | G. ’t Hooft, “Rigorous construction of planar diagram field theories in four-dimensional Euclidean space,” Commun. Math. Phys. **88** (1983) 1–25. |
| [Tri57]   | F. G. Tricomi, *Integral equations*, Interscience, New York (1957). |
| [VG88]    | J. C. Várilly and J. M. Gracia-Bondía, “Algebras of distributions suitable for phase space quantum mechanics. II: Topologies on the Moyal algebra,” J. Math. Phys. **29** (1988) 880–887. |
| [Wid38]   | D. V. Widder, “The Stieltjes transform,” Trans. Amer. Math. Soc. **43** (1938) 7–60. |
| [Wig56]   | A. S. Wightman, “Quantum field theory in terms of vacuum expectation values,” Phys. Rev. **101** (1956) 860–866. |
| [WK74]    | K. G. Wilson and J. B. Kogut, “The renormalization group and the $\epsilon$-expansion,” Phys. Rept. **12** (1974) 75–200. |
| [Wit91]   | E. Witten, “Two-dimensional gravity and intersection theory on moduli space,” Surveys Diff. Geom. **1** (1991) 243–310. |