Four Dimensional CFT Models with Rational Correlation Functions

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Abstract

Recently established rationality of correlation functions in a globally conformal invariant quantum field theory satisfying Wightman axioms is used to construct a family of soluble models in 4-dimensional Minkowski space-time. We consider in detail a model of a neutral scalar field $\phi$ of dimension 2. It depends on a positive real parameter $c$, an analogue of the Virasoro central charge, and admits for all (finite) $c$ an infinite number of conserved symmetric tensor currents. The operator product algebra of $\phi$ is shown to coincide with a simpler one, generated by a bilocal scalar field $V(x_1, x_2)$ of dimension $(1, 1)$. The modes of $V$ together with the unit operator span an infinite dimensional Lie algebra $\mathfrak{L}_V$ whose vacuum (i.e. zero energy lowest weight) representations only depend on the central charge $c$. Wightman positivity (i.e. unitarity of the representations of $\mathfrak{L}_V$) is proven to be equivalent to $c \in \mathbb{N}$.

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1 Introduction

The task of constructing a conformally invariant quantum field theory model - using dressed vertices and (global) operator product expansions (OPE) - has been set forth over 30 years ago ([26] [22] [25] [29] [30] [19] [34] [11] [35] [20] [16] [27] [31] [9] [17]; for a review of this early work and further references - see [37]). After a relative quiet (during which only some sporadic applications of the formalism appeared - see e.g. [7]) the subject has been gradually revived (see [36] [6] [32] [15] [24] [12] [10] among others) in the wake of the 2–dimensional (2D) conformal field theory (CFT) revolution (now the subject of textbooks– see e.g. [8] where a bibliography of original work can be found). It gathered new momentum with the discovery of the AdS–CFT correspondence and the associated intensified study of the $N = 4$ supersymmetric Yang-Mills theory (for a sample of recent papers and further references - see [1] [4]).

The present work is chiefly motivated by the concept of a rational conformal field theory (RCFT). Albeit this notion arose in the framework of 2D CFT, recent work [23] suggests that it may be relevant to any number of space–time dimensions. We consider in detail the simplest example beyond free fields, given in [23], the case of a model of a neutral scalar field of dimension 2. More complicated (and potentially more interesting) cases involving fields of dimension 3 and 4 are only briefly discussed.

We start by recalling the relevant results of [23] which allow us to derive the general expressions for the 4–point Wightman functions.

Adding to the Wightman axioms a condition of global conformal invariance (GCI) of local observables (i.e. invariance of correlation functions under a single–valued action of the 4–fold cover $G = SU(2, 2)$ of the conformal group whenever $x$ and $gx$ ($g \in G$) both belong to Minkowski space) we deduce the Huygens’ principle: local fields $\phi(x), \psi(y)$ commute whenever the difference $x - y$ is non–isotropic; moreover,

$$[(x - y)^2]^{-N/2} [\phi(x), \psi(y)] = 0 \quad \text{for} \quad N \gg 1 \quad \text{(1.1)}$$

(see [23] Theorem 4.1 and Proposition 4.3, where the precise bound for $N$ is given). This result is based on the fact that a space–like separated pair of points in Minkowski space can be mapped by a proper conformal transformation into a time-like one. (Thus, GCI is a stronger requirement than invariance of Schwinger functions under the euclidean conformal group.) The Huygens’ principle implies (together with energy positivity) that the Wightman distributions are rational functions of the form

$$W(x_1, \ldots, x_n) \equiv \langle 1, \ldots, n \rangle = P(x_1, \ldots, x_n) \prod_{1 \leq j < k \leq n} (\rho_{jk})^{-\mu_{jk}}, \quad \text{(1.2)}$$

where $P$ is a polynomial (in general, tensor valued),

$$x_{jk} \equiv x_j - x_k, \quad \rho_{jk} = x_{jk}^2 + i0x_{jk}^0 \quad \text{(1.3)}$$

([23] Theorem 3.1; the $i0x_{jk}^0$ is only essential when $\rho_{jk}$ occurs in denominators and prescribes the contour integration for Wightman distributions that reflects energy positivity - see [33]). Hilbert space positivity is taken into account using OPE and the classification of positive energy unitary irreducible representations of $G$ ([18]).

Expanding the discussion of Sec. 5 of [23] we shall derive the general form of the truncated 4–point function $W^4_4(d)$ of a neutral scalar field $\phi$ of integer dimension $d$ satisfying GCI (see Eq. (2.1) below).
Combining Proposition 5.3 and Corollary 4.4 of [23] we can write

\[ \mathcal{W}_d^i (d) \equiv \mathcal{W}^t (x_1, \ldots, x_4; d) = \mathcal{D}_d (\rho_{ij}) \mathcal{P}_d (\eta_1, \eta_2) , \]

\[ \mathcal{D}_d (\rho_{ij}) = \frac{(\rho_{13} \rho_{24})^{d-2}}{(\rho_{12} \rho_{23} \rho_{34} \rho_{14})^{d-1}} , \quad \mathcal{P}_d (\eta_1, \eta_2) = \sum_{i,j \geq 0} c_{ij} \eta_1^i \eta_2^j , \quad (1.4) \]

where \( \eta_i \) are the conformally invariant cross ratios

\[ \eta_1 = \frac{\rho_{12} \rho_{34}}{\rho_{13} \rho_{24}} , \quad \eta_2 = \frac{\rho_{14} \rho_{23}}{\rho_{13} \rho_{24}} . \quad (1.5) \]

For \( x_{jk}^2 \neq 0 \) we can ignore the \( i0x_{jk}^0 \) term in the definition of \( \rho_{jk} \) ([14]). The Huygens’ principle (strong locality) then implies symmetry under the permutation group \( S_4 \). Its normal subgroup \( \mathbb{Z}_2 \times \mathbb{Z}_2 \) (with non-trivial elements \( s_{12} s_{34}, s_{14} s_{23}, s_{13} s_{24} \), where \( s_{ij} \) is a substitution exchanging \( i \) and \( j \)) acts trivially on \( \eta_1 \) and \( \eta_2 \). Hence, it suffices to impose invariance under the \( 6 \)-element factor group \( S_4 / \mathbb{Z}_2 \times \mathbb{Z}_2 \cong S_3 \) generated by

\[ s_{12} : \mathcal{P}_d (\eta_1, \eta_2) \mapsto \eta_2^{d-3} \mathcal{P}_d \left( \frac{\eta_1}{\eta_2}, \frac{1}{\eta_2} \right) = \mathcal{P}_d (\eta_1, \eta_2) , \]

\[ s_{23} : \mathcal{P}_d (\eta_1, \eta_2) \mapsto \eta_1^{d-3} \mathcal{P}_d \left( \frac{1}{\eta_1}, \frac{\eta_2}{\eta_1} \right) = \mathcal{P}_d (\eta_1, \eta_2) \quad (1.6) \]

(which also involves \( s_{13} = s_{12} s_{23} s_{12} = s_{23} s_{12} s_{23} \) implying \( \mathcal{P}_d (\eta_2, \eta_1) = \mathcal{P}_d (\eta_1, \eta_2) \)). This leaves us with the following\[ \left\lfloor \frac{d^2}{3} \right\rfloor \] independent coefficients:

\[ c_{ij} \quad \text{for} \quad i \leq j \leq \frac{2d-3-i}{2} , \quad (1.7) \]

\[ (c_{ij} = c_{ij} = c_{i,2d-3-i-j} = c_{2d-3-i-j,i} = c_{j,2d-3-i-j} = c_{2d-3-i-j,j}) . \]

The present paper is chiefly devoted to the case \( d = 2 \), that is the minimal \( d \) for which a non–zero truncated 4–point function \( \mathcal{W}_d^i (d) \) exists. We shall set in this case\[ \left\lceil \frac{d^2}{3} \right\rceil \]

\[ \langle 12 \rangle = \frac{c_2}{2} (12)^2 , \quad \langle 123 \rangle = c_3 (12) (23) (13) , \]

\[ \mathcal{W}_d^i (d = 2) = c_4 (12) (23) (34) (14) (1 + \eta_1 + \eta_2) , \]

\[ (ij) = (4 \pi^2 \rho_{ij})^{-1} . \quad (1.8) \]

Parameters like

\[ c := \frac{c_3^2}{c_3^3} = 8 \frac{\langle 12 \rangle \langle 23 \rangle \langle 13 \rangle}{(\langle 123 \rangle)^2} , \quad c' := \frac{c_2}{c_4} \quad (1.9) \]

are invariant under rescaling of \( \phi \). It will be proven in Sec. 2 that if there is a single field \( (\phi) \) of dimension 2 then these constants are equal. Moreover, their common value \( c ( = c') \) also determines the normalization of the 2–point function of the stress–energy tensor and thus

1 \([a]\) stands for the integer part of \( a \) \( \left[ \frac{d^2}{3} \right] = 1, 3, 5, 8 \), for \( d = 2, 3, 4, 5 \) : \( \left[ \frac{(d+1)^2}{3} \right] - \left[ \frac{d^2}{3} \right] = \left[ \frac{d^2}{3} \right] (d+1) \) \( = 1, 2, 2, 3 \) for \( d = 1, 2, 3, 4 \).

2 The 4–point Wightman function obtained from (1.8) coincides with the one given by Proposition 5.3 and Eq. (5.16) of [24] for \( N_2 = \frac{c_2}{\pi^2}, C_2 = \frac{c_4}{(2 \pi)^2} \), \( C_{20} = C_{21} = 0 \).
appears as a generalization of the Virasoro central charge. We will then restrict our attention
to the case of a single field $\phi$ corresponding to $c_2 = c_3 = c_4 = c$.

Similarly, the general truncated 4–point function for $d = 3$ is

$$W'_4 (3) = \frac{\rho_{13} \rho_{24}}{(\rho_{12} \rho_{23} \rho_{34} \rho_{14})^2} \left\{ c_0 \left( 1 + \eta_1^2 + \eta_2^2 \right) + c_1 \left( (\eta_1 + \eta_2)(1 + \eta_1 \eta_2) + \eta_1^2 + \eta_2^2 \right) + b \eta_1 \eta_2 \right\}$$

$$= \{ c_i \equiv c_{0i} \text{ for } i = 0, 1 \text{, and } b \equiv c_{11} \}. \quad (1.10)$$

The requirement that no $d = 2$ (scalar) field is present in the OPE of two $\phi$'s in this case gives $c_1 = -c_0 \neq 0$, should one demand the presence of a stress energy tensor in the OPE.

The case $d = 4$ appears to be particularly interesting and will be briefly discussed in the concluding Sec. 6.

The paper is organized as follows.

In Sec. 2 we write down the OPE of two $\phi$'s in terms of a bilocal scalar field $V (x_1, x_2)$ of dimension $(1, 1)$ which satisfies – in each argument – the (free) d’Alembert equation. Using this result we sketch a proof of the statement that $V$ belongs to the OPE algebra generated by $\phi$, a property only valid in four space–time dimensions. The free field equations for $V$ then imply that the truncated $n$–point function of $\phi$ is expressed as a sum of 1–loop diagrams with propagators $(ij)$ and a common factor $c_n$ for all $n \geq 4$. The uniqueness of the field $\phi$ of dimension 2 is proven to correspond to $c_n = c \alpha^n$.

In Sec. 3 we establish the existence of an infinite set of conservation laws: the term with light cone singularity (12) (34) is reproduced by the contribution of an infinite number of (even rank) conserved symmetric traceless tensor currents

$$T_{2l} (x, \zeta) = T_{\mu_1 .. \mu_2l} (x) \zeta^{\mu_1} ... \zeta^{\mu_{2l}}, \quad \Box \zeta T_{2l} (x, \zeta) = 0 = \frac{\partial^2}{\partial x_\mu \partial \zeta^\mu} T_{2l} (x, \zeta), \quad (1.11)$$

to the OPE of two $\phi$'s (including the $l = 0$ term $T_0 (x) = \phi (x)$). For $\phi$ expressed as a linear combination of normal products of free fields

$$\phi (x) = \frac{1}{2} \sum_{i = 1}^N \alpha_i : \varphi_i^2 (x) : , \quad \langle 0 | \varphi_i (x_1) \varphi_j (x_2) | 0 \rangle = \delta_{ij} \quad (1.12)$$

the stress–energy tensor is also given by the sum of free field expressions:

$$T_2 (x, \zeta) = \sum_{i = 1}^N : \left\{ (\zeta \cdot \partial \varphi_i (x))^2 - \frac{1}{2} \zeta^2 \partial_\mu \varphi_i \partial^\mu \varphi_i + \frac{1}{6} [\zeta^2 \Box - (\zeta \cdot \partial)^2] \varphi_i^2 (x) \right\} : . \quad (1.13)$$

The case of equal $c_n = c$ ($n = 2, 3 ...$)– i.e. of a unique $\phi$ – corresponds to $\alpha_i = 1$ (for $i = 1, ..., N$) and $c = N$. The truncated $n$–point functions of $T_2 (x, \zeta)$ remain proportional to its free massless scalar field expression for all $c > 0$. Thus, the parameter $c$ indeed plays the role of a 4–dimensional extension of the Virasoro central charge.

In Sec. 4 we study the mode expansion of the bilocal field $V$ which naturally appears in the so called analytic compact picture. We exhibit an infinite dimensional Lie algebra $\mathfrak{L}_V$ spanned by the modes $V_{nm} (z_1, z_2)$ of $V$ and by the unit operator.

In Sec. 5 we prove that the unitary positive energy representations of $\mathfrak{L}_V$ correspond to positive integer $c$ (Theorem 5.1). Combining this theorem with Propositions 2.2 and 2.3 we
derive the same result for the original field algebra of the \( d = 2 \) scalar field \( \phi \). This implies that \( \phi \) belongs to the Borchers’ class of a system of free fields \(^{[3]}\) (see \(^{[33]}\) for a text-book introduction to this concept).

Sec. 6 is devoted to a discussion of the results. We indicate on the way how the methods of this paper apply to fields of dimension 3 and 4, and end up with the formulation of two open problems.

2 One loop \( n \)-point functions. OPE in terms of a bilocal field

We begin by rewriting the expression for the general 4–point function of a neutral scalar field \( \phi (x) \) of dimension 2 satisfying GCI in a form that suggests its generalization to the \( n \)-point function. According to (1.8) we have

\[
\langle 1234 \rangle = \langle 12 \rangle \langle 34 \rangle + \langle 13 \rangle \langle 24 \rangle + \langle 14 \rangle \langle 23 \rangle + W'_4 , \quad (\langle ij \rangle = \frac{c_2}{2} (ij)^2 ) ,
\]

where the truncated 4–point Wightman function can be written as a sum of contributions of three box diagrams:

\[
W'_4 = c_4 \{ (12) (34) (23) (14) + (12) (34) (13) (24) + (13) (24) (14) (23) \} .
\]

This expression is reproduced by an OPE for the product of two \( \phi \)'s that can be written compactly in terms of bilocal fields:

\[
\langle 0 | \phi (x_1) \phi (x_2) \rangle = \langle 0 | \{ (12) + (12) V (x_1, x_2) + : \phi (x_1) \phi (x_2) : \} \} , \quad V (x_1, x_2) = V (x_2, x_1) ,
\]

where the three terms are mutually orthogonal

\[
\langle 0 | V (x_1, x_2) | 0 \rangle = 0 = \langle 0 | : \phi (x_1) \phi (x_2) : | 0 \rangle = \langle 0 | V (x_1, x_2) : \phi (x_3) \phi (x_4) : | 0 \rangle ,
\]

and satisfy

\[
\langle 0 | V (x_1, x_2) V (x_3, x_4) | 0 \rangle = c_4 \{ (13) (24) + (14) (23) \} , \quad V (x_1, x_2) \phi (x_3) | 0 \rangle = c_3 (13) (23) , \quad : \phi (x_1) \phi (x_2) : \phi (x_3) \phi (x_4) | 0 \rangle = \langle 13 | 24 \rangle + \langle 14 | 23 \rangle + c_4 (13) (23) (14) (24) .
\]

(In general, a field \( V (x_1, x_2) \) is said to be bilocal if \([V (x_1, x_2), V (x_3, x_4)] = 0 \) for \( x_i \) space–like to \( x_j \), \( i = 1, 2, j = 3, 4 \) and if it commutes with all local fields \( \phi (x_3) \) of the theory for space–like \( x_{i3}, i = 1, 2 \).)

A priori, the algebra of \( V \) and \( : \phi (x_1) \phi (x_2) : \) may be larger than the OPE algebra of \( \phi \). It is a non–trivial result, valid only in 4–dimensions, that the (symmetric) bilocal fields \( V (x_1, x_2) \) and \( : \phi (x_1) \phi (x_2) : \) can actually be determined separately from the expansion (2.3).

**Proposition 2.1** If \( V (x_1, x_2) \) is a bilocal field obeying (2.5) then it satisfies in each argument the d’Alembert equation:

\[
\Box_1 V (x_1, x_2) = 0 = \Box_2 V (x_1, x_2) , \quad \Box_i = \frac{\partial^2}{\partial x_i^\mu \partial x_{i\mu}} , \quad i = 1, 2 ,
\]

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provided the metric in the state space is positive definite.

Proof. The vector valued distribution $\Box_i V(x_1, x_2)|0\rangle$, $i = 1, 2$, vanishes, due to Wightman positivity since the norm squares of the corresponding smeared vectors are expressed in terms of the 4–point function in the first equation \((2.3)\). The vanishing of $\Box_i V$ then follows from local commutativity by virtue of the Reeh–Schlieder theorem. (The argument is essentially the same as the proof of the statement that the vacuum is a separating vector for local fields– see [33] Sec. 4.) □

Proposition 2.2  The bilocal field

$W(x_1, x_2) := 4\pi^2 x_{12}^2 \{\phi(x_1)\phi(x_2) - \langle 12 \rangle\} = V(x_1, x_2) + 4\pi^2 x_{12}^2 :\phi(x_1)\phi(x_2):$ \((2.8)\)

allows to determine the Taylor coefficients in $x_1$ at $x_1 = x_2$ of the two terms in the right hand side separately.

Sketch of proof\(^3\) The (pseudo)harmonicity of $V$ \((2.5)\) implies

\[
(y\partial_1)^n W(x_1, x) \bigg|_{x_1=x} = (y\partial_1)^n V(x_1, x) \bigg|_{x_1=x} + y^2 n(n-1) 4\pi^2 :[(y\partial_x)^{n-2}\phi(x)]\phi(x):. \]

\((2.9)\)

In view of \((2.7)\) $\Box y(y\partial_1)^n V(x_1, x) \bigg|_{x_1=x} = 0$; thus $(y\partial_1)^n V(x_1, x) \bigg|_{x_1=x}$ appears as the harmonic part of the left hand side of \((2.9)\) viewed as a polynomial in $y$ and hence is uniquely determined by $W(x+y, x). □$

It is clear from \((2.3)\) that $V(x_1, x_2)$ is nonsingular for coinciding arguments. We can thus define a second local field

$\phi_2(x) = \frac{1}{2} V(x, x), \quad (2.10)$

of dimension 2; it can be a multiple of $\phi(x)$ only if the ratios \((1.9)\) coincide. Indeed, it follows from \((1.8)\) and \((2.3)\) that

$\langle 0 | \phi(x_1)\phi(x_2)|0\rangle = \langle 12 \rangle, \quad \langle 0 | \phi_2(x_1)\phi_2(x_2)|0\rangle = \frac{c_3}{c_2} \langle 12 \rangle, \quad \langle 0 | \phi_2(x_1)\phi_2(x_2)|0\rangle = \frac{c_4}{c_2} \langle 12 \rangle$;

thus

$\phi_2(x) = \lambda \phi(x) \left(= \frac{c_3}{c_2} \phi(x)\right) \text{ implies } c_2 c_4 = c_3^2. \quad (2.11)$

The preceding discussion admits an extension to the $n$–point truncated function. If we set, generalizing \((2.3)\),

$W_n^t(x_1, \ldots, x_n) = \frac{c_n}{2} \sum_{\sigma \in \text{Perm}(2...n)} (1\sigma_2) \underbrace{\sigma_2\sigma_3 \ldots \sigma_{n-1}\sigma_n}_{\sigma_{i\neq j}} (1\sigma_n),$

\(\sigma_{ij} = \begin{cases} (\sigma_i\sigma_j) & \text{for } \sigma_i < \sigma_j \\ (\sigma_j\sigma_i) & \text{for } \sigma_j < \sigma_i \end{cases}, \quad n = 2, 3, 4, \ldots . \quad (2.12)\)

then the field $\phi(x)$ of dimension 2 is unique if $c_n = c \alpha^n$ for some $\alpha > 0$, $n = 2, 3, \ldots$.

\(^3\)Complete proofs of Propositions 2.2 and 2.3 will be published elsewhere.
If we define $V_1$ as a linear combination of normal products of free (massless) fields,

$$V_1(x_1, x_2) = \sum_{i=1}^{N} \alpha_i :\varphi_i(x_1) \varphi_i(x_2):,$$  \hspace{1cm} (2.13)

and set $\phi(x) = \phi_1(x) = \frac{1}{2} V_1(x, x)$, then we can reproduce (2.12) with

$$c_n = \sum_{i=1}^{N} \alpha_i^n.$$  \hspace{1cm} (2.14)

Furthermore, we can introduce inductively a series of bilocal and local fields $V_n(x_1, x_2)$ and $\phi_n(x)$ of dimension $(1, 1)$ and 2 setting

$$V_n(x_1, x_2) = \lim_{x_{34} \to 0} \left\{ \frac{4 \pi^2 x_{34}^2}{\rho_{13}} \left[ V_1(x_1, x_3) V_{n-1}(x_2, x_4) - c_n \left( (12) (34) + (14) (32) \right) \right] \right\} = \sum_{i=1}^{N} \alpha_i^n :\varphi_i(x_1) \varphi_i(x_2):, \quad \phi_n(x) = \frac{1}{2} V_n(x, x).$$  \hspace{1cm} (2.15)

Note that the limit (2.15) is independent of the point $x_3 = x_4$ and that the field $V$ appearing in the OPE (2.3) coincides with $V_2$.

The dimension of the space of different $d = 2$ fields $\phi_k(x)$ is equal to the number of different values of $\alpha_i$ in (2.13). To see this we note that the Gram determinant of inner products

$$\langle 0 | \phi_j(x_1) \phi_k(x_2) | 0 \rangle = \frac{1}{2} (12)^2 \sum_{i=1}^{N} \alpha_i^{j+k}$$  \hspace{1cm} (2.16)

is a multiple of $\prod_{i=1}^{N} \alpha_i^2 \prod_{1 \leq j < k \leq N} (\alpha_j - \alpha_k)^2$.

**Remark 2.1** Fields of type (2.13) have been studied in a different context (for bounded 2–dimensional fields) in [3] [13] where also infinite sums are admitted. We restrict our discussion to finite $N$ since only in this case a stress energy tensor exists – and is given by (1.13).

From now on we shall restrict our discussion to the simplest case of a single field $\phi$ of dimension 2 and set

$$c_n = c \quad \text{for} \quad n = 2, 3, 4, ...$$  \hspace{1cm} (2.17)

(absorbing the possible factor $\alpha^n$ in the normalization of $\phi$).

The general form (2.12) of the truncated $n$–point function can in fact, be deduced.

**Proposition 2.3** Let $\phi(x)$ be a GCI Wightman field of dimension 2 whose truncated $n$–point function is given by (2.12) with $c_n = c$ for $n \leq 4$. Then the limit

$$V(x_1, x_2) = \lim_{\rho_{13} \to 0} \rho_{13} \rho_{23} \left\{ \phi(x_1) \phi(x_2) \phi(x_3) - \langle 13 \rangle \phi(x_2) - \langle 23 \rangle \phi(x_1) - \langle 123 \rangle \right\}$$  \hspace{1cm} (2.18)

exists, does not depend on $x_3$, and defines a harmonic in each argument bilocal field $V(x_1, x_2)$. Furthermore, the truncated $n$–point functions of $\phi$ will be given by (2.12) for all $n$. 7
Sketch of proof. Equation (2.2) and the conservation of the stress-energy tensor (see Sec. 3) imply that (2.12) is valid for \( n \leq 6 \). The 1-loop expression for the 6-point function allows to derive (2.18). The expression for the correlation function (2.5) of two \( V \)'s satisfies the d’Alembert equation in each argument. By virtue of Proposition 2.1 the operator field \( V \) obeys this equation in its entire domain. Eq. (2.12) for \( n \leq 6 \) implies an expansion of the form

\[
\phi(x_1) \phi(x_2) \phi(x_3)|0\rangle = \langle 123 |0\rangle + \sum_{i=1,2,3} \left\{ \langle jk | \phi(x_i) + \sum_{i \neq j, j \neq k} i j \ k V(x_j, x_k) + (jk) :V(x_j, x_k) \phi(x_i): \right\} |0\rangle +
\]

\[
: \phi(x_1) \phi(x_2) \phi(x_3): |0\rangle
\]

((i, j, k) form permutations of 1, 2, 3). The result then follows. □

Remark 2.2 If we drop the requirement of Wightman positivity - which implies the validity of the stress-energy tensor conservation as an operator equation - then the general form of the truncated 5-point function would be

\[
\mathcal{W}_5^t(x_1, ..., x_5) = \lambda \mathcal{W}_5^t(2.12)+4 \pi^2 c (1 - \lambda) \sum_{1 \leq i < j \leq 5} \rho_{ij} \prod_{1 \leq k \leq 5, j \neq k} k j, \ \lambda \in \mathbb{R}.
\]

We note that the 1–dimensional time–like restriction \( \phi(t, 0) \) of \( \phi(x) \) satisfies all properties of the chiral stress energy tensor in a 2D CFT. It follows that all restricted truncated functions should have the form (2.12). This is satisfied by (2.20) (for our choice of constants) because of a non–trivial identity between the two terms in the 1–dimensional case.

Corollary 2.4 Under the assumptions of Proposition 2.3 one can prove (using also Proposition 2.2) that the field algebra of \( \phi(x) \) coincides with the algebra of the bilocal field \( V(x_1, x_2) \).

Demanding that the truncated \( n \)–point function of \( \phi \) for \( n \geq 3 \) is strictly less singular in \( x_{ij} \) than its 2–point function we have taken into account a necessary condition for Wightman positivity. We shall prove a necessary and sufficient condition for positivity in Sec. 5.

Remark 2.3 If we rescale the field \( \phi \) by a factor \( c^{-\frac{1}{2}} \) and let \( c \rightarrow \infty \) we recover the case of a generalized free field of dimension 2:

if \( \hat{\phi}(x) = \frac{1}{\sqrt{c}} \phi(x) \) then

\[
\lim_{c \rightarrow \infty} \langle 0 | \hat{\phi}(x_1) \hat{\phi}(x_2) \hat{\phi}(x_3) \hat{\phi}(x_4) |0\rangle = \langle 12 \rangle_1 \langle 34 \rangle_1 + \langle 13 \rangle_1 \langle 24 \rangle_1 + \langle 14 \rangle_1 \langle 23 \rangle_1,
\]

where \( \langle ij \rangle_1 = \frac{1}{2} (ij)^2 \).

3 Expansion of \( V(x_1, x_2) \) in local fields. Infinite set of conserved tensor currents

We shall now demonstrate that our model possesses an infinite number of conserved local tensor currents. More precisely, the bilocal field \( V(x_1, x_2) \) can be expanded in a series of even rank,
conserved symmetric traceless tensor fields $T_{2l} (x, \zeta)$ (1.11) (of twist $= \text{dimension} - \text{rank} = 2$):

$$V (x_1, x_2) = 2 \sum_{l=0}^{\infty} C_l K_l (x_{12} \cdot \partial_2 , \rho_{12} \Box_2) T_{2l} (x_2, x_{12}) , \quad (3.1)$$

reproducing the 4–point function (2.5). Here

$$K_l (s, t) = \frac{(2l + 1)!}{(l)!^2} \int_0^1 d\alpha \, (1 - \alpha)^l e^{\alpha s} \sum_{n=0}^{\infty} \frac{(-\frac{\alpha (1-\alpha)}{4})^n}{n! (2l + 1)^n} , \quad (K_l (0, 0) = 1) , \quad (3.2)$$

$\partial_2$ is the derivative in $x_2$ for fixed $x_{12}$, $\Box_2$ is the corresponding d’Alembert operator, $(\nu)_n = \frac{\Gamma(n + \nu)}{\Gamma(\nu)}$; it is chosen to transform the 2–point function $\langle 0 | T_{2l} (x_2, \zeta_2) T_{2l} (x_3, \zeta_3) | 0 \rangle$ into a 3–point function:

$$K_l (x_{12} \cdot \partial_2, \rho_{12} \Box_2) \frac{\rho_{23}^{2l+2}}{\rho_1 \rho_2} (x_{12} \cdot r (x_{23}) \cdot \zeta)^{2l} = \frac{(X \cdot \zeta)^{2l}}{\rho_1 \rho_2} , \quad (3.3)$$

where

$$\zeta \cdot r (x_{23}) \cdot \zeta = \zeta \cdot \zeta - \frac{\langle \zeta \cdot x_{23} \rangle \langle \zeta \cdot x_{23} \rangle}{\rho_2} , \quad X := X_{12}^3 := \frac{x_{13}}{\rho_{13}} - \frac{x_{23}}{\rho_{23}} , \quad (3.4)$$

In verifying (3.3) (see [10]) one applies the relation

$$\left( \frac{\Box_2}{4} \right)_n (y \cdot \zeta)^m = \frac{\Gamma(n + m - 1)}{\Gamma(n + \nu)} (y \cdot \zeta)^m \quad \text{for} \quad \zeta^2 = 0$$

(used for $y = x_{23} + \alpha x_{12}$). In order to compute individual contribution of $T_{2l}$ to the 4–point function of $\phi$ we need the 3–point function

$$\langle 0 | \phi (x_1) \phi (x_2) T_{2l} (x_3, \zeta) | 0 \rangle = N_l C_l (12) (X^2)^{l+1} (\zeta^2)^l C_{2l}^1 (\hat{X} \cdot \hat{\zeta}) , \quad \hat{X} := \frac{X}{\sqrt{X^2}} , \quad (3.5)$$

where $N_l > 0$, $C_{n}^1 (z)$ is the Gegenbauer polynomial satisfying

$$\left\{ (1 - z^2) \frac{d^2}{dz^2} - 3z \frac{d}{dz} + n (n + 2) \right\} C_{n}^1 (z) = 0 , \quad C_{n}^1 (1) = n + 1 . \quad (3.6)$$

Writing the normalization constant in (3.3) as a product, $N_l C_l$, we exploit the fact that the 3–point function vanishes whenever the structure constant $C_l = 0$.

**Remark 3.1** The homogeneous polynomial $H_{2l} (x, \zeta) = (x^2 \zeta^2)^l C_{2l}^1 (\hat{X} \cdot \hat{\zeta})$ is the harmonic extension of the monomial $(2x \cdot \zeta)^{2l}$ defined on the light cone $\zeta^2 = 0$ (cf. [2]):

$$\Box_2 H_{2l} (x, \zeta) = (x^2 \zeta^2)^l (\zeta^2)^{l-1} \times \left\{ (1 - z^2) \frac{d^2}{dz^2} C_{2l}^1 (z) - 3z \frac{d}{dz} C_{2l}^1 (z) + 4l (l + 1) C_{2l}^1 (z) \right\} = 0$$

(for $z = \hat{X} \cdot \hat{\zeta}$), \quad $H_{2l} (x, \zeta) \bigg|_{\zeta^2 = 0} = (2x \cdot \zeta)^{2l} . \quad (3.7)$
Similarly, the 2-point function $\langle 0| T_{2l}(x_1, \zeta_1) T_{2l}(x_2, \zeta_2)| 0 \rangle$ is proportional to $\rho_{12}^{-2l-2} (\zeta_1^2 \zeta_2^2)^l C_{2l}^1 \left( \frac{\zeta_1 \cdot \zeta_2 - 2(\zeta_1 \cdot x_{12})(\zeta_2 \cdot x_{12})}{\rho_{12}} \right)$.

Inserting (2.3) in the 4-point function (2.1) (2.2) and using (2.5) and the expansion (3.1) for $V(x_3, x_4)$ we find

$$\langle 0| \phi(x_1) \phi(x_2) V(x_3, x_4)| 0 \rangle = c \langle 12 \rangle ((13)(24) + (14)(23)) =$$

$$= 2 \sum_{l=0}^{\infty} C_l K_l(x_{34} \cdot \partial_{4}, \rho_{34} \Box_{4}) \langle 0| \phi(x_1) \phi(x_2) T_{2l}(x_4, x_{34})| 0 \rangle =$$

$$= 4 \langle 12 \rangle \sum_{l=0}^{\infty} N_l C_l^2 \left( \frac{4l+1)!}{(2l)!^2} \int_0^1 \alpha 2^l (1 - \alpha)^{2l} \left( -\frac{\alpha (1 - \alpha)}{4} \rho_{34} \Box_{4} \right)^n \right. \times$$

$$\times \rho_{34}^l \left( X_y^2 \right)^{l+1} C_{2l}^1 \left( \tilde{X}_y \cdot \tilde{x}_{34} \right) ,$$

$$X_y = \frac{x_1 - y}{\rho_{1y}} - \frac{x_2 - y}{\rho_{2y}} , \quad y = x_4 + \alpha x_{34} ,$$

$$\rho_{iy} = \rho_{i4}(1 - \alpha) + \alpha \rho_{i3} - \alpha (1 - \alpha) \rho_{34} \quad , \quad i = 1, 2 .$$

(3.8)

It will be convenient for what follows to substitute the second conformally invariant cross ratio $\eta_2 (1.3)$ by the difference $\epsilon = 1 - \eta_2$ which tends to zero for $x_{34} \rightarrow 0$ (or $x_{12} \rightarrow 0$):

$$\epsilon = 1 - \eta_2 ( = O(x_{34}) = O(x_{12}) ) .$$

(3.9)

**Proposition 3.1** For

$$N_l C_l^2 = \frac{(4l)!}{2l!}^{-1}$$

(3.10)

the contribution of $V(x_3, x_4)$ to the 4-point function (2.1) is reproduced by the superposition (3.8) of 3-point functions of the twist 2 fields $T_{2l}$

$$\frac{\langle 0| V(x_1, x_2) V(x_3, x_4)| 0 \rangle}{(13)(24)} = c \left( 1 + \frac{1}{1 - \epsilon} \right) =$$

$$= 2c \sum_{l=0}^{\infty} (4l + 1) \int_0^1 \left[ \frac{\epsilon \alpha (1 - \alpha)}{1 - \epsilon \alpha} \right]^{2l} d\alpha \frac{1}{1 - \epsilon \alpha} .$$

(3.11)

The proof of this statement is given in Appendix A.

The Ward–Takahashi identity for the time–ordered 3–point function of the stress–energy tensor allows to compute the normalization $N_1 C_1$ of the Wightman function (3.3):

$$\langle 0| \phi(x_1) \phi(x_2) T_2(x_3, \zeta)| 0 \rangle = \frac{(12)}{3 \pi^2} X^2 \left( X^2 \zeta^2 - 4 (X \cdot \zeta)^2 \right) , \quad \text{i.e.} \quad N_1 C_1 = \frac{-1}{3 \pi^2} .$$

(3.12)
Comparing with (3.10) we find that both $N_1$ and $C_1$ are transcendental, only the combination $N_1C_1^2$ being rational:

$$C_1 = -\frac{\pi^2}{2}, \quad N_1 = \frac{2}{3\pi^4}, \quad (N_1C_1^2 = \frac{1}{6}). \quad (3.13)$$

**Remark 3.2** It is instructive to note that the contribution of each $T_{2l}$ to the ratio (3.11) (given by the $l$th term in the right hand side) involves a logarithmic function in $1 - \epsilon$ (see Appendix A) while the infinite sum is a rational function of $\epsilon$.

## 4 The infinite dimensional Lie algebra of field modes and its bilocal realization

The conformal compactification $\overline{M} = S^3 \times S^1 / \mathbb{Z}_2$ of Minkowski space $M = \mathbb{R}^{3,1}$ gives rise to a natural notion of conformal energy, the generator of (isometric) rotation of the time-like circle $S^1$, and of an associated discrete basis of field modes. We shall parametrize $\overline{M}$ following (3.11) in terms of complex coordinates $z = (z_a, a = 1, 2, 3, 4)$ fixed by the involution $z \mapsto z^* := \frac{\overline{z}}{z^2}$:

$$\overline{M} = \left\{ z = (z_a \in \mathbb{C}, a = 1, ..., 4) ; \quad z^*_a := \frac{\overline{z}_a}{z^2} = z_a \quad (z^2 = \sum_a z^*_a =: z^2 + z_4^2) \right\}. \quad (4.1)$$

This condition implies the property

$$z^2 \overline{z}^2 = 1, \quad \frac{z_a z_b}{z^2} = \overline{z}_a z_b = z_a \overline{z}_b \in \mathbb{R} \quad \text{for} \quad z \in \overline{M} \quad (4.2)$$

which, in turn, characterizes this parametrization of $\overline{M}$. We choose the embedding map $M \subset \overline{M}$ as

$$M \ni (x^0, x) \mapsto z = \omega^{-1}(x) x, \quad z_4 = \frac{1-x^2}{2\omega(x)}, \quad \omega(x) = \frac{1+x^2}{2} - ix^0. \quad (4.3)$$

Clearly, $z$ defined by (4.3) satisfies (4.2); in particular,

$$z^2 = \frac{\omega(x)}{\omega(x)} = \frac{(1+ix^0)^2 + x^2}{(1-ix^0)^2 + x^2} = \frac{1}{\overline{z}^2}, \quad |z^2| = z \cdot \overline{z} = 1 \quad \text{(for} \quad z \in \overline{M}). \quad (4.4)$$

In order to write down the inverse transformation it is convenient to present $z$ in terms of a complex quaternion (or, equivalently, an element of $U(2)$ – see [38]):

$$q z = z_4 + z q, \quad q_i q_j = \epsilon_{ijk} q_k - \delta_{ij} \quad \text{(i.e.} \quad q_1 q_2 = -q_2 q_1 = q_3, \quad \text{etc.}) \quad (4.5)$$

The cone at infinity, $K_\infty = \overline{M} \setminus M$, consists of the quaternions $qz \in \overline{M}$ for which $1 + qz$ is not invertible

$$q z \in K_\infty \iff 2\omega^{-1} := (1 + q z) (1 + q^+ z) = (1 + z_4)^2 + z^2 = 0 \quad (q^+ z = z_4 - q z). \quad (4.6)$$
For \( qz \notin K_\infty \) we can set
\[
ix : = ix_0 + qx = \frac{qz - 1}{qz + 1} \quad \text{or} \quad ix_0 = \omega_z \frac{z^2 - 1}{2}, \quad x = \omega_z z = \frac{2z}{(1 + z_4)^2 + z^2}.
\] (4.7)

We shall use the fact that the flat metric on \( \overline{M} \) is related to the Poincaré invariant metric on \( M \) by the complex conformal factor \( \omega \) [4.3]
\[
dz^2 = dz^2 + dz_4^2 = \omega^{-2}(x) \, dx^2 \quad (dx^2 = dx^2 - dx_4^2).
\] (4.8)

To a scalar field \( \phi_M(x) \) of dimension \( d \) in Minkowski space we make correspond an analytic \( z \)-picture field \( \phi(z) \) defined by:
\[
\phi(z) = (2\pi)^d \omega_z^d \phi_M(x(z)) \quad (\omega_z = \frac{2}{(1 + z_4)^2 + z^2} = \omega(x(z))).
\] (4.9)

for \( x(z) \) given by (4.7). The term analytic is justified by the fact that energy positivity implies analyticity of the vector valued function \( \phi(z)|0 \) for \( |z|^2 < 1 \). Indeed, the future tube \( T_+ = \{ \xi \in \mathbb{C}^4; \text{Im} \xi^0 > |\text{Im} \xi| \} \), the analyticity domain of \( \phi_M(\xi)|0 \) (see [33]), is mapped into a complex neighbourhood \( \Sigma_+ \) of the 4-dimensional unit ball \( \mathbb{B}^4 \); more precisely, we have
\[
\mathbb{B}^4 = \{ \xi \in \mathbb{R}^4; \xi^2 := \xi_0^2 + \xi_4^2 < 1 \}, \quad \mathbb{B}^4 \times S^1 / \mathbb{Z}_2 = \{ \xi \in \mathbb{B}^4, \tau \in \mathbb{R} \} \subset \Sigma_+.
\] (4.10)

Note that \( \overline{M} \) appears as the boundary of the 5-dimensional manifold \( \mathbb{B}^4 \times S^1 / \mathbb{Z}_2 \):
\[
z \in \overline{M} \quad \text{iff} \quad z = e^{\tau} \hat{z}, \quad \tau \in \mathbb{R}, \quad \hat{z} \in S^3 = \{ \hat{z} \in \mathbb{R}^4; \hat{z}^2 = 1 \}. \quad (4.11)
\]

The conformal Hamiltonian \( H \) is, in this picture, nothing but the (hermitian) generator of translation in \( \tau \):
\[
e^{iHt} \phi(z) = e^{-iHt} \phi(e^{it}z) \quad \text{or} \quad [H, \phi(z)] = \left( d + z_a \frac{\partial}{\partial z_a} \right) \phi(z), \quad H|0 \rangle = 0.
\] (4.12)

The decomposition of \( \phi \) into eigenmodes of \( H \) reads
\[
\phi(z) = \sum_{n \in \mathbb{Z}} \phi_n(z), \quad [\phi_n(z), H] = n \phi_n(z).
\] (4.13)

The modes \( \phi_n(z) \) can be written as power series in \( z_a \) and \( \frac{1}{z_4} \) that are homogeneous in \( z \) of degree \( -n - d \).

For a free field \( \varphi(z) \) of dimension \( d = 1 \) the modes \( \varphi_{\pm n} \) are homogeneous harmonic polynomials spanning a space of dimension \( n^2 \) (as a space of \( \text{SO}(4) \) symmetric traceless tensors of rank \( n - 1 \): \( \binom{n + 2}{3} - \binom{n}{3} = n^2 \)); in particular, \( \varphi_0(z) = 0, \varphi_1(z) = \varphi_{\pm 1}(z) = a_{\pm 1}, \varphi_2(z) = \varphi_{\pm 2}(z) = a_{\pm 2} z_\mu \). They are subject to the canonical commutation relations [36]
\[
[\varphi_n(z), \varphi_m(w)] = \frac{(w^2)^{n-1}}{(z^2)^{n+1}} C_{|n|-1}^{1} \hat{z} \cdot \hat{w} \, \delta_{n,-m}, \quad (z = \sqrt{z^2})
\]
\[
\epsilon(n) = \begin{cases} 
1 & \text{for } n > 0 \\
0 & \text{for } n = 0 \\
-1 & \text{for } n < 0 
\end{cases}.
\] (4.14)
we can write down a basis in the vacuum Hilbert space $\phi$ of the analyticity of the vector valued function $V$. The OPE follows from (4.10) and (4.15) that $n$ and $\{23\}$, which, as a result, cannot obey a free field equation. It is all the more remarkable to note that the 2–point function $\langle 0 | \phi(z) \phi(w) | 0 \rangle$ appears as a generating function for the Gegenbauer polynomials defined in (3.6).

One can expand the bilocal field $V$ in modes $V = \sum_{n\neq 0} V_{nm}$, which behave as products of $\phi$–modes:

$$
\Delta_z V_{nm}(z, w) = 0 = \Delta_w V_{nm}(z, w),
$$

$$
\left( z \frac{\partial}{\partial z} + n + 1 \right) V_{nm}(z, w) = 0 = \left( w \frac{\partial}{\partial w} + m + 1 \right) V_{nm}(z, w).
$$

(4.15)

(The homogeneity condition only agrees with the Laplace equation if we set $V_{0m} = 0 = V_{nm}$.) The modes of the $d = 2$ field $\phi$ are most conveniently expressed as infinite sums of $V$–modes:

$$
2 \phi_n(z) = \sum_{\nu \in \mathbb{Z}} V_{\nu, n-\nu}(z, z) \quad (V_{mn}(z, z) = V_{nm}(z, z))
$$

(4.16)

The components $V_{\nu, n-\nu}(z, z)$ of $\phi_n(z)$ (unlike those of $\phi_n(z)$) span an infinite dimensional space. This is a common feature for scalar fields of dimension $d > 1$ (more generally, for elementary conformal fields of weight $(j_1, j_2; d)$ with $d \geq j_1 + j_2 + 2$, in the notation of [13] and [23], which, as a result, cannot obey a free field equation). It is all the more remarkable that the state space for a given energy $n$ is always finite dimensional. This is a consequence of the analyticity of the vector valued function $V_{nm}(z, w)|0\rangle$ for $z, w \in \mathbb{C}_\pm$. Indeed, it then follows from (4.10) and (1.13) that

$$
V_{nm}(z, w)|0\rangle = 0 \quad \text{if} \quad n \geq 0 \quad \text{or} \quad m \geq 0.
$$

(4.17)

Consequently, only $(n-1)$ terms of the infinite sum (4.16) contribute to the vector $\phi_{-n}(z)|0\rangle$:

$$
2 \phi_{-n}(z)|0\rangle = \sum_{\nu=1}^{n-1} V_{\nu, n-\nu}(z, z)|0\rangle.
$$

In order to display the identity of the vacuum state spaces of $\phi$ and $V$, guaranteed by Corollary 2.4 we need to include the composite twist 2 fields $T_{2l}(z, \zeta)$ in the operator algebra of $\phi$. Here is the realization of the four lowest energy spaces in the two pictures. Setting for the vacuum Hilbert space

$$
\mathcal{H} = \mathcal{H}_0 \oplus \bigoplus_{n=2}^{\infty} \mathcal{H}_n, \quad (H - n) \mathcal{H}_n = 0 \quad (\text{dim} \mathcal{H}_0 = \text{dim} \mathcal{H}_2 = 1)
$$

we can write down a basis in $\mathcal{H}_2, \mathcal{H}_3$ and $\mathcal{H}_4$ as follows:

$$
\phi_{-2}|0\rangle = \frac{1}{2} V_{-1,-1}|0\rangle; \quad \phi_{-3}^a z_a|0\rangle = V_{-2,-1}^a z_a|0\rangle \quad (z_a V_{-1,-2}^a|0\rangle); \quad \{ \phi_{-4}^{ab}|0\rangle, T_{2l}(0, \zeta)|0\rangle, \phi_{-2}^2|0\rangle \} \sim \{ V_{-2,-2}^{ab} z_a z_b|0\rangle, V_{-3,-1}^{ab} z_a z_b|0\rangle, V_{-1,-1}^2|0\rangle \}.
$$

The difficulty in describing the full state space $\mathcal{H}$ in such a manner stems from the fact that the modes of $\phi$ do not span an (infinite dimensional) Lie algebra: the commutator $[\phi(z_1), \phi(z_2)]$ also involves all twist 2 conserved tensors $T_{2l}(z_2, z_{12})$ (and their derivatives in the first argument). $T_{2l}$ ($l = 0, 1, ...$) together with the unit operator exhaust, in fact, the singular terms in the OPE $\phi(z) \phi(w)|0\rangle$. The resulting commutator algebra simplifies drastically for collinear $z_j = \zeta_j e (e^2 = 1)$: it then reduces to the Virasoro algebra:

$$
[L_n, L_m] = (n - m) L_{n+m} + \frac{c}{12} n (n^2 - 1) \delta_{n,-m} \quad \text{for} \quad \phi_n(\zeta e) = \frac{L_n}{\zeta^{n+2}}, \quad L_n = \phi_n(e).
$$

(4.18)
The point is that the second argument, \( z_{12} \), of \( T_2 \) cancels the singular factor \( \frac{1}{z_{12}} \) in the OPE in the 1–dimensional case.

Using the orthogonality of different quasipinary fields we can produce a sample of projected commutation relations between \( \phi_n (z) \) for non–collinear arguments illustrating the appearance of the Virasoro subalgebra as a special case.

To begin with we note that the vacuum OPE \((2.3)\) remains valid in the \( z \)–picture provided we set

\[
(12) = \frac{1}{z_{12}^2} \quad \text{(implying} \quad \langle 12 \rangle = \frac{c}{2} (12)^2 = \frac{c}{2} (z_{12}^2)^{-2}, \quad z^2 = z^2 + z_4^2 \quad \text{)}
\]

(4.19)

(the singularity at \( z_{12}^2 = 0 \) being treated as a limit from the domain \( |z_1^2| > |z_2^2| \)). Using the knowledge of the generating function for the Gegenbauer polynomials,

\[
\left( \frac{1}{(z-w)^2} \right) ^\lambda = \frac{1}{(z^2)^\lambda} \left( 1 - 2 \hat{\phi} \cdot \hat{w} \sqrt{\frac{w^2}{z^2} + \frac{z^2}{w^2}} \right) ^{-\lambda} = \frac{1}{(z^2)^\lambda} \sum_{n=0}^\infty \left( \frac{w^2}{z^2} \right)^n C_n^\lambda (\hat{\phi} \cdot \hat{w}) ,
\]

and the expressions \((1.8)\) and \((2.2)\) for \( 2–, 3–, \) and \( 4– \) point correlation functions of \( \phi \) we can write the term involving the central extension of the Lie algebra generated by \( \phi_n \):

\[
\langle 0 | \phi_2 [ \phi_n (z), \phi_{-n} (w) ] \phi_{-2} | 0 \rangle = \frac{(w^2)^{\hat{\phi} - 1}}{(z^2)^{\hat{w}}} \times
\]

\[
\times \langle 0 | \phi_2 \left\{ C_{n-2}^1 (\hat{\phi} \cdot \hat{w}) \phi_0 (z) + \frac{w^2}{z^2} C_n^1 (\hat{\phi} \cdot \hat{w}) \phi_0 (w) + \frac{c}{2 \hat{\phi} \cdot \hat{w}} C_{n-2}^2 (\hat{\phi} \cdot \hat{w}) \right\} \phi_{-2} | 0 \rangle
\]

\[
n \geq 1, \quad (C_{-1}^\lambda \equiv 0). \quad (4.21)
\]

The Virasoro subalgebra \((4.18)\) is recovered for collinear arguments noting the normalization property for Gegenbauer polynomials:

\[
C_n^\lambda (1) = \left( n + 2 \lambda - 1 \right) = \frac{(2 \lambda)_n}{n!} , \quad \left( \frac{2}{n} + \frac{\nu}{\nu+1} \right) \sum_{\nu=1} \nu C_{\nu-l}^l (1) = 2 \mu + 2 - \nu \quad .
\]

(In particular, Eq. \((4.21)\)) reproduces \((4.18)\) for \( n + m = 0 \).

The Lie algebra \( \Sigma_V \) of the bilocal field \( V \) is much simpler to describe. The modes \( V_{nm} \) of \( V \) satisfying \((1.15)\) and the unit operator span by themselves an infinite dimensional Lie algebra:

\[
[V_{n_1 n_2} (z_1, z_2), V_{n_3 n_4} (z_3, z_4)] =
\]

\[
= c \prod_{j=1}^4 \left( z_j^2 \right) ^{-n_j + 1} \left\{ C_{[n_1] - 1}^1 (\hat{z}_1 \cdot \hat{z}_2) C_{[n_2] - 1}^1 (\hat{z}_2 \cdot \hat{z}_4) \delta_{n_1, -n_3} \delta_{n_2, -n_4} +
\right.
\]

\[
+ C_{[n_1] - 1}^1 (\hat{z}_1 \cdot \hat{z}_3) C_{[n_2] - 1}^1 (\hat{z}_2 \cdot \hat{z}_4) \delta_{n_1, -n_4} \delta_{n_2, -n_3} \} \epsilon (n_1) \epsilon (n_2) +
\]

\[
+ (z_1^2)^{-n_1 + 1} (z_3^2)^{-n_1 + 1} C_{[n_1] - 1}^1 (\hat{z}_1 \cdot \hat{z}_3) \epsilon (n_1) \delta_{n_1, -n_2} V_{n_2 n_4} (z_2, z_4) +
\]

\[
+ (z_2^2)^{-n_2 + 1} (z_3^2)^{-n_2 + 1} C_{[n_2] - 1}^1 (\hat{z}_2 \cdot \hat{z}_3) \epsilon (n_2) \delta_{n_2, -n_3} V_{n_1 n_4} (z_1, z_4) +
\]

\[
+ (z_1^2)^{-n_1 + 1} (z_4^2)^{-n_1 + 1} C_{[n_1] - 1}^1 (\hat{z}_1 \cdot \hat{z}_4) \epsilon (n_1) \delta_{n_1, -n_3} V_{n_1 n_2} (z_1, z_3) +
\]

\[
+ (z_2^2)^{-n_2 + 1} (z_4^2)^{-n_2 + 1} C_{[n_2] - 1}^1 (\hat{z}_2 \cdot \hat{z}_4) \epsilon (n_2) \delta_{n_2, -n_4} V_{n_1 n_3} (z_1, z_3) .
\]

(4.23)
It is, in fact, a central extension of the infinite dimensional real symplectic algebra \( sp(\infty, \mathbb{R}) \). According to (4.16) the \( \phi \)-modes belong to this algebra. The vacuum representation of \( \mathfrak{L}_V \) is characterized by the energy positivity condition (4.17).

The associative algebra of \( V_{nm}(z,w) \) contains an ideal \( \mathcal{I}_0 \) generated by

\[
\{ V_{n0}(z,w) \ (= V_{0n}(w,z)) ; \ n \in \mathbb{Z} \} \ (\in \mathcal{I}_0) .
\]

which annihilates all states in the vector space \( \mathcal{H}_V \) spanned by polynomials in \( V_{-n,-m} \ (n, m \in \mathbb{N}) \) acting on the vacuum. Although \( \mathcal{I}_0 \) may well be represented non-trivially in other sectors of the theory it is natural to work with the factor algebra \( \mathcal{B}_V \) in the vacuum sector. Indeed, \( \mathcal{B}_V \) can be identified as the operator algebra, generated by the bilocal field \( \phi \) in \( B \) of the theory acting (non-trivially) on \( \mathcal{H}_V \). The relative simplicity of the operator algebra \( \mathcal{B}_V \) in \( \mathcal{H}_V \) stems from the fact that the modes \( V_{nm}(z,w) \ (n \neq 0 \neq m) \) are (homogeneous) harmonic functions in \( z \) and \( w \) see (4.15). It follows from our analysis of the mode space of the free field \( \varphi(z) \) that \( V_{nm}(z,w) \) span a space of dimension \( n^2 m^2 \) except for the diagonal, \( n = m \), for which the symmetry of \( V \) implies that the dimension of the space is \( \binom{n^2 + 1}{2} \).

The modes \( V_{nm} \) are eigenvectors of the Cartan elements

\[
h_l = \frac{l}{2 \pi^2} \int V_{-l,l}(u,u) \delta \left( \sqrt{u^2 - 1} \right) \, d^4u , \quad (u^2 = u^2 + u_4^2) , \quad l \in \mathbb{N} . \quad (4.25)
\]

(Parametrizing \( u \in \mathbb{S}^3 \) by \( u = (\sin \psi \sin \theta \cos \varphi, \sin \psi \sin \theta, \sin \psi \cos \theta, \cos \psi) \) we can replace the volume element \( \delta \left( \sqrt{u^2 - 1} \right) \, d^4u \) by \( \sin^2 \psi \sin \theta \, d\psi \, d\theta \, d\varphi \), \( 0 \leq \psi \leq \pi \), \( 0 \leq \theta \leq \pi \), \( 0 \leq \varphi \leq 2\pi \); the normalization factor \( \frac{1}{2\pi^2} \) fixes the integral (of 1) over \( \mathbb{S}^3 \) to 1.) We have, in particular,

\[
(h_l - \delta_{lm} - \delta_{ln}) \ V_{-n,-m}(\tilde{z},\tilde{w}) | 0 \rangle = 0 \quad (\text{for } n, m \in \mathbb{N}) . \quad (4.26)
\]

In deriving this property one uses the relation

\[
\frac{l}{2 \pi^2} \int C_{l-1}^1(\tilde{w} \cdot u) \ C_{n-1}^1(\tilde{u} \cdot \tilde{z}) \delta \left( \sqrt{u^2 - 1} \right) \, d^4u = \delta_{ln} C_{n-1}^1(\tilde{w} \cdot \tilde{z}) . \quad (4.27)
\]

It follows that the conformal Hamiltonian \( H \) defined in (4.12) can be written in the form

\[
H = \sum_{l=1}^{\infty} l \ h_l . \quad (4.28)
\]

## 5 Unitary vacuum representations of \( \mathfrak{L}_V \)

We begin by introducing an antiinvolution in \( \mathcal{B}_V \) and the associated inner product in \( \mathcal{H} \).

We define a star operator in the algebra of modes setting

\[
V_{nm}(z,w)^* = V_{-m,-n}(w,z) \ (= V_{-n,-m}(z,w)) \quad \text{for} \quad z, w \in \mathbb{M} , \quad (5.1)
\]

so that \( V(z,w)^* = V(z,w) \).

**Remark 5.1** The antiinvolution (5.1) involves a correspondence between homogeneous harmonic functions of degree \( n - 1 \) and \( -n - 1 \). If we write, for \( n, m > 0 \),

\[
V_{-n,-m}(z,w) = V_{-n,-m}^{b_1 \ldots b_{n-1}, a_1 \ldots a_{m-1}} z_{b_1} \ldots z_{b_{n-1}} w_{a_1} \ldots w_{a_{m-1}}
\]
then we shall have

\[ V_{-n,-m}(z, w)^* = V_{nm}(w, z) = \frac{1}{w^2 z^2} V_{mn}^{a_1 \ldots a_{m-1}, b_1 \ldots b_{n-1}} \frac{w_{a_1}}{w^2} \ldots \frac{w_{a_{m-1}}}{w^2} \frac{z_{b_1}}{z^2} \ldots \frac{z_{b_{n-1}}}{z^2}, \]

where both \( V_{-n,-m} \) and \( V_{nm} \) are symmetric traceless tensors of rank \((n-1, m-1)\) (with respect to the indices \(a_i\) and \(b_j\), separately).

We shall call a Hilbert space \((\mathcal{H})\) representation of \(\mathfrak{L}_V\) unitary if the (positive) scalar product in \(\mathcal{H}\) and the conjugation \((\cdot)\) in \(\mathfrak{L}_V\) are related by

\((\Phi, X \Psi) = (X^* \Phi, \Psi)\) for every \(X \in \mathfrak{L}_V\), \(\Phi, \Psi \in \mathcal{H}^F\),

\[ (5.2) \]

where \(\mathcal{H}^F\) is the dense subspace of finite energy vectors of \(\mathcal{H}\) which belongs to the domain of any \(X\) in \(\mathfrak{L}_V\).

One can introduce a (not necessarily positive) inner product \(\langle \cdot, \cdot \rangle\) in \(\mathcal{H}_V\) satisfying \((5.2)\) defining the bra vacuum by conditions conjugate to \((4.17)\):

\[ \langle 0 | V_{nm} = 0 \text{ unless } n > 0 \text{ and } m > 0, \]

\[ (5.3) \]

and assuming \(\langle 0 | 0 \rangle = 1\). The main result of this section is the following characterization of the unitary vacuum representation of \(\mathfrak{L}_V\).

**Theorem 5.1** The inner product in \(\mathcal{H}_V\), defined for a (normalized) vacuum vector satisfying \((1.17)\) and \((5.3)\) and for \(V_{nm}(z, w)\) obeying \((4.23)\), is positive semidefinite iff \(c \in \mathbb{Z}_+ = \{0, 1, 2, \ldots\}\).

**Proof.** Fix a unit vector \(e \in \mathbb{S}^3\) and consider the 1–dimensional subalgebra \(\mathfrak{L}_V^e\) of \(\mathfrak{L}_V\) generated by

\[ v_{nm} := V_{nm}(e, e) \in \mathfrak{L}_V^e \subset \mathfrak{L}_V, \quad n, m \in \mathbb{Z}, \quad e^2 = 1. \]

\[ (5.4) \]

It follows from \((4.23)\) and from \((4.22)\) that \(v_{nm}\) satisfy the commutation relations of the modes of a 1–dimensional (chiral) bilocal current:

\[ [v_{n_1 m_1}, v_{n_2 m_2}] = c n_1 m_1 (\delta_{n_1, -n_2} \delta_{m_1, -m_2} + \delta_{n_1, -m_2} \delta_{m_1, -n_2}) + \]

\[ + n_1 (\delta_{n_1, -n_2} v_{m_1 m_2} + \delta_{n_1, -m_2} v_{m_1 n_2}) + \]

\[ + m_1 (\delta_{m_1, -n_2} v_{n_1 m_2} + \delta_{m_1, -m_2} v_{n_1 n_2}). \]

\[ (5.5) \]

**Lemma 5.2** There is a vector \(|\Delta_n\rangle \in \mathcal{H}_V^{(n(n+1))}\) whose norm square is a multiple of \(c (c-1) \ldots (c-n+1)\):

\[ \langle \Delta_n | \Delta_n \rangle \equiv ||| \Delta_n \||^2 = (n+1)! c (c-1) \ldots (c-n+1). \]

\[ (5.6) \]
Proof. It follows from (5.5) that the norm square of a polynomial of degree \( n \) in \( v_{kl} \) is a polynomial of degree (not exceeding) \( n \) in \( c \). We shall demonstrate that \( \langle \Delta_n | \Delta_n \rangle \) vanishes for integer \( c \) in the interval \( 0 \leq c < n \). To this end we note that if \( c \) is a positive integer and \( \mathcal{J}_m \), \( m \in \mathbb{Z} \) are \( c \)-dimensional operator valued vectors \( \mathcal{J}_m = \{ J^i_m, \ i = 1, ..., c \} \) satisfying
\[
\left[ J^i_m, J^j_n \right] = m \delta_{m,-n} \delta_{ij}, \quad m, n \in \mathbb{Z}, \quad i, j = 1, ..., c,
\] (5.7)
then the normal products
\[
v^{(c)}_{lm} = : \mathcal{J}_l : \mathcal{J}_m : = \sum_{i=1}^{c} : J^i_l J^i_m :
\] (5.8)
satisfy the commutation relations (5.5). If \( c < n \) then \( \det (v_{ij}) \) appearing in the definition of \( \langle \Delta_n \rangle \), which is the Gram determinant of the scalar products of \( n \) vectors in a \( c \)-dimensional space, should vanish. The coefficient \((n + 1)!\) to the leading \((n)\) power of \( c \) is computed as a sum of norm squares of terms entering the expansion of the determinant; for instance, for \( n = 4 \) we have
\[
\lim_{c \to \infty} \left( \frac{1}{c!} \langle \Delta_4 | \Delta_4 \rangle \right) = \frac{1}{4!^2 c^4} \left\{ \| \langle 0 | V_{11} \ldots V_{44} \|^2 + 6 \| \langle 0 | V^2_{12} V_{33} V_{44} \|^2 + \right.
\]
\[
\quad + 4 \| \langle 0 | V_{12} V_{23} V_{34} V_{44} \|^2 + 3 \| \langle 0 | V^2_{12} V^2_{34} \|^2 + \right.
\]
\[
\quad + 3 \| 2 \langle 0 | V_{12} V_{23} V_{34} V_{44} \|^2 \right\} =
\]
\[
= 2^4 + 6 \times 8 + 4 \times 8 + 3 \times 4 + 3 \times 4 = 120 (= 5!) .
\]

Remark 5.2 The Lie algebra \( \mathcal{L}_\nu \) of bilocal modes, characterized by the commutation relations (4.23) has a reductive star subalgebra \( \mathcal{U}_\infty \) (with no central extension) generated by \( V_{-n,m} (z, w) \), \( n, m \in \mathbb{N} \):
\[
[V_{-n_1,m_1} (\hat{z}_1, \hat{w}_1), V_{-n_2,m_2} (\hat{z}_2, \hat{w}_2)] =
\]
\[
= C^1_{m_1-1} (\hat{w}_1 \cdot \hat{z}_2) \delta_{m_1 n_2} V_{-n_1,m_2} (\hat{z}_1, \hat{w}_2) - C^1_{m_2-1} (\hat{w}_2 \cdot \hat{z}_1) \delta_{m_2 n_1} V_{-n_2,m_1} (\hat{z}_2, \hat{w}_1) ,
\] (5.9)
with a central element
\[
C_1 = \sum_{n=1}^{\infty} h_n ,
\] (5.10)
where \( h_n \) are the Cartan operators (4.25). We have
\[
[V_{-l,m} (\hat{z}, \hat{w}), C_1] =
\]
\[
= \frac{1}{2 \pi^2} \int \left\{ m C^1_{m-1} (\hat{w} \cdot u) V_{-l,m} (\hat{z}, u) - \right.
\]
\[
- l C^1_{l-1} (\hat{z} \cdot u) V_{-l,m} (u, \hat{w}) \right\} \delta \left( \sqrt{u^2} - 1 \right) d^4 u = 0 ,
\] (5.11)
where we again used the relation \((4.27)\). \(U_\infty\) contains what could be called the Cartan subalgebra of \(L_V\) spanned by the elements \(V_{-n,n} (e, e)\) for \(n \in \mathbb{N}\), \(e^2 = 1\) (including \(h_l\) \((4.25)\)). \(L_V\) is compounded by \(U_\infty\), the unit element and by a pair of conjugate abelian subalgerbas \(L^\pm\) (which are \(U_\infty\) modules), spanned by

\[
L^+ \supset \{ V_{-n,-m} (z, w) \}, \quad L^- \supset \{ V_{nm} (z, w) \}, \quad n, m \in \mathbb{N}.
\]

\(L^+\) consists of positive, \(L^-\), of negative root vectors with respect to the Cartan elements \(h_l\) \((4.25)\):

\[
[h_l, V_{\mp n, \mp m} (\hat{z}, \hat{w})] = \pm (\delta_{in} + \delta_{im}) V_{\pm n, \pm m} (\hat{z}, \hat{w}) .
\]

The commutators between elements of \(L^-\) and \(L^+\) belong to \(U_\infty \cup c1\). The operator

\[
C_2 = \sum_{n,m=1}^{\infty} \frac{n m}{8 \pi^4} \int\int V_{-n,-m} (v, u) \, V_{nm} (u, v) \, \delta \left( \sqrt{u^2} - 1 \right) \delta \left( \sqrt{v^2} - 1 \right) \, d^4u \, d^4v
\]

commutes with \(U_\infty\) and should have a positive spectrum in any unitary representation of \(L_V\). The counterpart of \(C_2\) \((5.14)\) for the subalgebra \(L^c\),

\[
C^c_2 = \frac{1}{2} \sum_{n,m \geq 1} \frac{1}{n m} \, v_{-n,-m} v_{nm}
\]

has its minimal eigenvalue in the subspace

\[
\mathcal{H}^{(n)}_c = \{ P_n (v_{-k,-l}) | 0 \rangle ; \ P_n \text{ homogeneous of degree } n \text{ in } v_{-k,-l} \}
\]
on the vector \(| \Delta_n \rangle\) (conjugate to) \((5.6)\):

\[
C^c_2 | \Delta_n \rangle = n (c - n + 1) | \Delta_n \rangle , \quad [C^c_2 - n (c - n + 1)] \big|_{\mathcal{H}^{(n)}_c} \geq 0 .
\]

We have, for instance,

\[
(C^c_2 - n [c + 2 (n - 1)]) \, v_{-k,-k}^n | 0 \rangle = 0 = (C^c_2 - n c) \, v_{-2n,-(2n-1)} \cdots v_{-2,-1} | 0 \rangle .
\]

It follows from Lemma 5.2 that there exist negative norm vectors unless \(c\) is a positive integer. To prove that for \(c \in \mathbb{N}\) the vacuum representation of \(L_V\) is indeed unitary it suffices to note that in this case \(V\) can be written in the form

\[
V (z_1, z_2) = \sum_{i=1}^{c} : \varphi_i (z_1) \varphi_i (z_2) :.
\]

where \(\varphi_i\) are mutually commuting free zero mass fields and to recall that a system of free fields satisfies all Wightman axioms (including positivity).

We have established on the way the following result (as a direct consequence of Lemma 5.2).

\[
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\]
Proposition 5.3 The vacuum representation of the infinite dimensional Lie algebra \( \mathfrak{L}_v \) of the 2–dimensional (2D) bilocal chiral field

\[
v(z, w) = \frac{1}{z w} \sum_{n, m \in \mathbb{Z}} v_{nm} z^{-n} w^{-m}, \quad z, w \in \mathbb{C}
\]

(5.18)

whose modes satisfy (5.3) (and \( v_{nm} |0\rangle = 0 \) unless \( m < 0 \) and \( n < 0 \)) is only unitary for positive integer \( c \).

This is an analogue of Kac–Radul theorem \([14]\) on the unitary representations of the \( W_{1+\infty} \) algebra. It is clear that the algebra of the 2D stress tensor

\[
T(z) = \frac{1}{2} v(z, z) = \sum_{n \in \mathbb{Z}} L_n z^{-n-2},
\]

(5.19)

– i.e. the Virasoro algebra \([4,18]\) – is a true subalgebra of \( \mathfrak{L}_v \) since it admits unitary representations for all \( c \geq 1 \) as well as a discrete series for \( c = c_n = 1 - \frac{6}{(n+1)(n+2)} \) \( (n = 1, 2, ...) \), the unitary Virasoro module \( \mathcal{H}_{c_n} \) being the quotient space of the corresponding lowest weight module with respect to a singular vector at "level" (= eigenvalue of \( L_0 \) \( n (n+1) \)).

The situation is different for \( D = 4 \) since \( V(z, w) \) is harmonic in each argument in that case. Due to Corollary 2.4 the algebra \( \mathcal{B}_V \) is then not bigger than the original OPE algebra of \( \phi \) so that the result of Theorem 5.1 extends to it.

Corollary 5.4 Under the assumptions of Proposition 2.3 it follows from Theorem 5.1 that the quantum theory of the field \( \phi \) with truncated \( n \)–point function (2.12) satisfies Wightman positivity iff \( c \) is a natural number (in which case \( \phi \) belongs to the Borchers’ class of a set of free fields).

6 Extensions of the results. Concluding remarks

The preceding results – and methods – apply to fields of higher dimension and arbitrary tensor structure. We shall establish important special cases of the following

Conjecture. If a neutral tensor field of integer dimension has truncated \( n \)–point functions which are multiples of the corresponding correlators of normal products of (derivatives of) free fields for \( n \leq 6 \), then Wightman positivity implies that the proportionality constant is a positive integer. (As indicated in Sec. 2, for \( d = 2 \) the statement follows from the expression for the 4–point function.)

Our first example is a conserved current whose (first five) truncated correlation functions are obtained from those of the current of a system of 2–component spinors,

\[
J^\mu (x; c_\psi) = \sum_{j=1}^{c_\psi} :\bar{\psi}_j^\dagger (x) \bar{\sigma}^\mu \psi_j (x):, \quad (\bar{\sigma}^0) = \bar{\sigma}_0 = 1 = \sigma_0, \quad \bar{\sigma}^j = -\sigma^j = -\sigma_j,
\]

(6.1)
by substituting the positive integer \( c_\psi \) by an arbitrary real number. Here \( \psi_j \) are mutually anticommuting free Weyl fields:

\[
\langle 0 | \psi_j (x_1) \psi_k^* (x_2) | 0 \rangle = \delta_{jk} S (x_{12}) , \quad S (x_{12}) = i \tilde{\partial}_2 (12) = i \frac{x_{12}}{2 \pi^2 \rho_{12}^3} , \tag{6.2}
\]

and we have used the conventions

\[
\tilde{\partial}_2 = \sigma_\mu \frac{\partial}{\partial x_{2 \mu}} \quad (x = \sigma_\mu x^\mu) , \quad \sigma^\mu \sigma_\nu + \sigma_\nu \sigma^\mu = -2 \delta^\mu_\nu . \tag{6.3}
\]

Introducing the spin–tensor components of the current

\[
J (x) (= J_{\alpha \beta} (x)) = \frac{1}{2} \sigma_\mu J^\mu (= \sum_{j=1}^{c_\psi} : \psi_j a (x) \psi_j^* (x) : ) . \tag{6.4}
\]

we can write

\[
\langle 0 | J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) | 0 \rangle = c_\psi S_{\alpha_1 \beta_1} (x_{12}) \ t S_{\beta_1 \alpha_2} (x_{12}) = c_\psi \left\{ S_{\alpha_1 \beta_1} (x_{12}) S_{\alpha_2 \beta_2} (x_{12}) - \frac{c_{\alpha_1} c_{\alpha_2} c_{\beta_1} c_{\beta_2}}{4 \pi^4 \rho_{12}^3} \right\} , \tag{6.5}
\]

\[
J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) - \langle 0 | J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) | 0 \rangle = t S_{\beta_1 \alpha_2} (x_{12}) V_{\alpha_1 \beta_2} (x_1, x_2) + S_{\alpha_1 \beta_2} (x_{12}) V_{\beta_1 \alpha_2} (x_1, x_2) + : J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) : , \tag{6.6}
\]

where \( ^t S \ ( ^t V \) stands for the transposed of \( S \ ( V \). Multiplying both sides by \( \frac{2 \pi^2}{i} \rho_{12} \tilde{x}_{12} \beta_{2 \alpha_2} \) and setting

\[
W_{\alpha_1 \beta_1} (x_1, x_2) = \frac{2 \pi^2}{i} \rho_{12} \tilde{x}_{12} \beta_{2 \alpha_2} \ \left\{ J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) - \langle 0 | J_{\alpha_1 \beta_1} (x_1) J_{\alpha_2 \beta_2} (x_2) | 0 \rangle \right\} \tag{6.7}
\]

we obtain

\[
W_{\alpha_1 \beta_1} (x_1, x_2) = V_{\alpha_1 \beta_1} (x_1, x_2) + ^t V_{\beta_1 \alpha_1} (x_1, x_2) + : J_{\alpha_1 \beta_1} (x_1) \text{ tr} (\tilde{x}_{12} J (x_2)) : \tag{6.8}
\]

where the bilocal field \( V \) satisfies

\[
\langle 0 | V_{\alpha_1 \beta_1} (x_1, x_2) V_{\alpha_2 \beta_2} (x_3, x_4) | 0 \rangle = S_{\alpha_1 \beta_2} (x_{14}) \ t S_{\beta_1 \alpha_2} (x_{23}) . \tag{6.9}
\]

It follows from (6.9) that

\[
\tilde{\partial}_1 \hat{\alpha}_{1 \alpha_1} V_{\alpha_1 \beta_1} (x_1, x_2) = 0 = V_{\alpha_1 \beta_1} (x_1, x_2) \ \frac{\partial}{\partial x_{2 \mu}} \ \tilde{\sigma} \hat{\beta}_1 \beta . \tag{6.10}
\]

As a result \( V_{\alpha \beta} \) and the normal product of \( J \) appearing in the right hand side of (6.9) can be determined separately and we can prove as in Sec. 5 that \( c_\psi (c_\psi - 1) \ldots (c_\psi - n + 1) \geq 0 \) for \( n = 1, 2, \ldots \).

As a second example we consider the Lagrangian density

\[
\mathcal{L}_F (x) = - \frac{1}{4} \sum_{\alpha=1}^{c_\psi} : F_{\mu \nu}^{\alpha} (x) F_{\alpha}^{\mu \nu} (x) : \quad ( c_F \in \mathbb{N} ) \tag{6.11}
\]

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and the associated analytic continuation of truncated Wightman functions to arbitrary positive real $c_F$. The truncated $n$–point function of $\mathcal{L}_F$ can again be written as a sum of $\frac{1}{2} (n - 1)!$ 1–loop graphs, the propagator associated with a line joining the vertices 1 and 2 being

$$D_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_{12}) = \frac{1}{4} \left\{ \partial_{\lambda_1} \partial_{\lambda_2} - \partial_{\mu_1} \partial_{\mu_2} - \partial_{\mu_1} \partial_{\lambda_2} - \partial_{\lambda_1} \partial_{\mu_2} \right\} \frac{1}{4 \pi^2 \rho_{12}} =$$

$$= \frac{r_{\lambda_1 \lambda_2} (x_{12}) r_{\mu_1 \mu_2} (x_{12}) - r_{\lambda_1 \mu_2} (x_{12}) r_{\mu_1 \lambda_2} (x_{12})}{4 \pi^2 \rho_{12}^2}. \quad (6.12)$$

This expression for the propagator also enters the OPE of two $\mathcal{L}$’s (together with a tensor valued bilocal field):

$$\langle 0| \mathcal{L}_F (x_1) \mathcal{L}_F (x_2) = \langle 0| \left\{ 2 c_F D_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_{12}) D_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_{12}) + \right.$$

$$\left. + D_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_{12}) V_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_1, x_2) + : \mathcal{L}_F (x_1) \mathcal{L}_F (x_2) : \right\},$$

$$2 D_{\lambda_1 \mu_1 \lambda_2 \mu_2} D_{\lambda_1 \mu_1 \lambda_2 \mu_2} = \frac{3}{(\pi \rho_{12})^4}. \quad (6.13)$$

For $c \in \mathbb{N}$, $V$ has a realization as a sum of normal products of free Maxwell fields:

$$V_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_1, x_2) = : F_{\lambda_1 \mu_1} (x_1) F_{\lambda_2 \mu_2} (x_2) : . \quad (6.14)$$

The OPE (6.13) allows to compute the truncated 4–point function of $\mathcal{L}_F$ which appears as a special case of the 5–parameter expression $\mathcal{W}_d^t (d = 4)$ computed from Eqs. (1.4)–(1.7):

$$\mathcal{W}_d^t (4) = \frac{\rho_{13}^2 \rho_{24}^2}{\rho_{12}^3 \rho_{23}^3 \rho_{34} \rho_{14}} \left\{ c_0 (1 + \eta_1^5 + \eta_2^5) + c_1 (\eta_1 + \eta_2 + \eta_1^4 + \eta_2^4 + \eta_1 \eta_2 (\eta_1^3 + \eta_2^3)) + \right.$$

$$\left. + c_2 (\eta_1^2 + \eta_2^2 + \eta_1^3 + \eta_2^3 + \eta_1 \eta_2 (\eta_1 + \eta_2)) + \right.$$

$$\left. + b_1 \eta_1 \eta_2 (1 + \eta_1^2 + \eta_2^2) + b_2 \eta_1 \eta_2 (\eta_1 + \eta_2) \right\},$$

$$(c_i \equiv c_{0i} \quad \text{for} \quad i = 0, 1, 2; \quad b_i \equiv c_{1i} \quad \text{for} \quad i = 1, 2). \quad (6.15)$$

Indeed the contribution $\mathcal{W}_d$ of the box diagram (computed by using formulae for traces of products of $r_{\mu_\nu}$ given in Appendix B),

$$\mathcal{W}_d = c_F D_{\lambda_1 \mu_1 \lambda_2 \mu_2} (x_{12}) D_{\lambda_1 \lambda_4} \mu_{14} (x_{14}) D_{\lambda_2 \lambda_3 \lambda_{14}} \mu_{23} (x_{23}) D_{\lambda_3 \lambda_4 \mu_3} (x_{34}) =$$

$$= 32 c_F \frac{(12) (23) (34) (14)}{\rho_{12} \rho_{23} \rho_{34} \rho_{14}} \left( 1 + \frac{\eta_1}{\eta_2} + \frac{\eta_2}{\eta_1} + \frac{\eta_1}{\eta_2} - \frac{2}{\eta_1} + \frac{2}{\eta_1 \eta_2} \right), \quad (6.16)$$

which enters the expression for the truncated 4–point function of $\mathcal{L} (x)$

$$\mathcal{W}_d^t (4) = (1 + s_{12} + s_{23}) \mathcal{W}_d (x_1, x_2, x_3, x_4) =$$

$$= \mathcal{W}_d (x_1, x_2, x_3, x_4) + \mathcal{W}_d (x_2, x_1, x_3, x_4) + \mathcal{W}_d (x_1, x_3, x_2, x_4) \quad (6.17)$$

fits the expression (6.15) for

$$c_0 = c_2 = b_1 = -\frac{1}{2} c_1 = \frac{c_F}{8 \pi^8}, \quad b_2 = 0. \quad (6.18)$$

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The first local field in the expansion of $V$ around the diagonal is the stress energy tensor:

$$T_{\nu} = \frac{1}{4} V^{\kappa\lambda}_{\kappa\lambda} (x, x) \delta_{\nu}^{\mu} - V^{\lambda\mu}_{\lambda\mu} (x, x) = - \mathcal{L} (x) \delta_{\nu}^{\mu} - V^{\lambda\mu}_{\lambda\mu} (x, x). \quad (6.19)$$

Conversely, the bilocal tensor field $V^{\lambda_{1}\mu_{1}}_{\lambda_{2}\mu_{2}} (x_{1}, x_{2}) = - V^{\mu_{1}\lambda_{1}}_{\lambda_{2}\mu_{2}} (x_{1}, x_{2}) = - V^{\lambda_{1}\mu_{1}}_{\mu_{2}\lambda_{2}} (x_{1}, x_{2})$ appears in the OPE of two $T_{\nu}^{\mu}$ and can be determined from it in two steps. First, one derives the formula

$$\langle 0 | V^{\lambda_{1}\mu_{1}}_{\lambda_{2}\mu_{2}} (x_{1}, x_{2}) V^{\lambda_{3}\mu_{3}}_{\lambda_{4}\mu_{4}} (x_{3}, x_{4}) | 0 \rangle = c_{F} \mathcal{D}^{\lambda_{1}\mu_{1}\lambda_{3}\mu_{3}} (x_{13}) \mathcal{D}_{\lambda_{2}\mu_{2}\lambda_{4}\mu_{4}} (x_{24}) + c_{F} \mathcal{D}^{\lambda_{1}\mu_{1}}_{\lambda_{4}\mu_{4}} (x_{14}) \mathcal{D}^{\lambda_{3}\mu_{3}}_{\lambda_{2}\mu_{2}} (x_{23}). \quad (6.20)$$

and deduces from it that $V^{\lambda_{1}\mu_{1}}_{\lambda_{2}\mu_{2}} (x_{1}, x_{2})$ satisfies in each argument the free Maxwell equations. Secondly, one uses this fact to single out the contribution of $V$ in the OPE of two $T$'s. Once more Wightman positivity implies $c_{F} \in \mathbb{N}$.

**Remark 6.1** The use of different notation, $c ( = c_{\phi})$, $c_{\psi}$ and $c_{F}$ for the constants multiplying the truncated functions of normal products of the free fields $\phi$, $\psi$ and $F_{\mu\nu}$, respectively, is justified by the fact that they correspond to (and exhaust the) different tensor structures in the general conformal invariant 3–point function of the stress energy tensor [32].

At the same time the 4–point functions of the conserved current $J_{\mu}$ and $\mathcal{L} (x)$ involve structures which cannot be reduced to normal products of free fields. If, for instance, $b_{2} \neq 0$ in (6.13) the 3–point function of $\mathcal{L} (x)$ won’t vanish (unlike the case of superposition of type (6.11) of normal products of free Maxwell fields). More generally, we have a 4–parameter family of admissible 4–point functions of $\mathcal{L} (x)$ obtained from Eq. (6.15) with the restriction

$$c_{2} = - c_{0} - c_{1} ( \neq 2 c_{0}) \quad (6.21)$$

coming from the requirement that no $d = 2$ field appears in the OPE of $\mathcal{L} (x_{1}) \mathcal{L} (x_{2})$ (and that the stress energy tensor is present in this OPE). They are only compatible with 3–point functions of $T$ of the type (6.19) (i.e. with the third of the three admissible structures in this 3–point function given in [32] cf. Remark 6.1).

To summarize: looking for a 4–dimensional RCFT beyond the Borchers’ class of free fields we have excluded the theory of a bilocal field of dimension (1, 1) and have come to the following problem. Assume that the only local fields in the observable algebra, satisfying GCI, of dimension $d \leq 4$ are the (conserved traceless) stress energy tensor $T_{\mu\nu} (x)$ and a scalar field $\mathcal{L} (x)$ of dimension 4 (playing the role of an action density). The problem is to construct an OPE algebra consistent with the $n$–point functions of these fields for $n \leq 4$ that would allow to compute higher point correlation functions and to implement the condition of Wightman positivity. This example is attractive because the dimensions of the basic fields $\mathcal{L}$ and $T_{\mu\nu}$ are protected. Moreover, in any renormalizable quantum field theory one can define a (gauge invariant) local action density and a stress energy tensor.
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Appendix A  Proof of Proposition 3.1

We shall first compute the sum in the right hand side of (3.8) for 
\[ \rho_{34} = 0, \quad 2 X_y \cdot x_{34} = \frac{\rho_{14} - \rho_{13}}{(1 - \alpha) \rho_{14} + \alpha \rho_{13}} = \frac{\rho_{24} - \rho_{23}}{(1 - \alpha) \rho_{24} + \alpha \rho_{23}} = \frac{\rho_{14} \rho_{23} - \rho_{13} \rho_{24}}{[(1 - \alpha) \rho_{14} + \alpha \rho_{13}][(1 - \alpha) \rho_{24} + \alpha \rho_{23}]} = \frac{1}{\rho_{13} \rho_{24}} X_y^2 = \frac{1}{\alpha + (1 - \alpha) \frac{\rho_{14}}{\rho_{13}}} [1 - \rho_{23} \rho_{24}] \cdot \frac{1}{\alpha + (1 - \alpha) \frac{\rho_{14}}{\rho_{13}}} [1 - \alpha + \frac{\rho_{23}}{\rho_{24}}], \]  
\[ \frac{\rho_{13} \rho_{24}}{\rho_{12}} X_y^2 = \frac{1}{\alpha + (1 - \alpha) \frac{\rho_{14}}{\rho_{13}}} [1 - \alpha + \frac{\rho_{23}}{\rho_{24}}], \]  
\[ (A.1) \]
and then use the result to give a general proof of Proposition 3.1. According to (A.1) we have
\[ \lim_{\rho_{34} \to 0} \left( \rho_{34} X_y^2 \right)^l C_{2l}^1 \left( \tilde{X}_y \cdot \tilde{x}_{34} \right) = \frac{\epsilon^{2l}}{\alpha + (1 - \alpha) \frac{\rho_{14}}{\rho_{13}}} [1 - \rho_{23} \rho_{24}] \frac{\epsilon^{2l}}{1 - \alpha + \frac{\rho_{23}}{\rho_{24}}}. \]
Conformal invariance allows to send \( x_1 \) to infinity setting \( \frac{\rho_{14}}{\rho_{13}} \to 1, \frac{\rho_{24}}{\rho_{23}} \to 1 - \epsilon \) thus reproducing the right hand side of (3.11). Taking the sum in \( l \) we reduce the proof of (3.11) to verifying the identity
\[ 2 \int_0^1 \frac{(1 - \epsilon \alpha) [(1 - \epsilon \alpha)^2 + \epsilon^2 \alpha^2 (1 - \alpha)^2]}{(1 - \epsilon \alpha^2) (1 - 2 \epsilon \alpha + \epsilon \alpha^2)} \, d\alpha = 1 + \frac{1}{1 - \epsilon} \]  
\[ (A.2) \]
which is straightforward.

It is also instructive to compute the individual terms in the right hand side of (3.11) which correspond to the contribution of twist 2 fields to the OPE. Using Euler’s integral representation for the hypergeometric function we find
\[ 1 + \frac{1}{1 - \epsilon} = 2 \sum_{l=0}^{\infty} \left( \frac{4l}{2l} \right)^{-1} e^{2l} F(2l + 1, 2l + 1; 4l + 2; \epsilon). \]  
\[ (A.3) \]
Each \( F(2l + 1, 2l + 1; 4l + 2; \epsilon) \) is, in fact, an elementary function. In particular, the first two terms which provide the contribution of the original field \( \phi \) and of the stress–energy tensor \( T_2 \) to the OPE can be written in the form
\[ 2 \int_0^1 \left. \frac{\langle 0 | V(x_1, x_2) \phi(x_4 + \alpha x_{34}) | 0 \rangle}{c(13)(24)} \, d\alpha \right|_{\rho_{34} = 0} = 2 F(1, 1; 2; \epsilon) = \]
\[
2 C_1 \left( \int_0^1 \frac{\mathcal{V}(x_1, x_2) T_2(x_1 + \alpha x_{34}, x_{34})}{e^{13} (24)} \, d\alpha \right) \right|_{\rho_{34} = 0} = \frac{\epsilon^2}{3} F(3, 3; 6; \epsilon) = \frac{60}{\epsilon^2} \left( \frac{1}{\epsilon} - 1 + \frac{\epsilon}{6} \ln \frac{1}{1 - \epsilon} - 1 + \frac{\epsilon}{2} \right) = \epsilon^2 \left\{ \frac{1}{3} + \frac{\epsilon}{2} + \sum_{n=2}^{\infty} \frac{(4)^l n^2}{(3)^n (7)^{n-1}} \right\}. 
\]

(A.4)

Proceeding to the general case \((\rho_{34} \neq 0)\) we shall use the following generalization of (A.3) (see [10]). Exchange the conformal cross ratios \((1.3) (3.9)\) with the variables \(\eta\) and \(\bar{\eta}\) related to \(\eta_1\) and \(\epsilon\) by

\[
\eta \bar{\eta} = \eta_1, \quad \eta + \bar{\eta} = \epsilon + \eta_1, \quad ((1 - \eta)(1 - \bar{\eta}) = \eta_2) .
\]

(A.5)

(We note that for space–like \(x_{ij}\) the variables \(\eta\) and \(\bar{\eta}\) are complex conjugate to each other.) In terms of these variables we can write (see Eq. (3.10) of [10]):

\[
\frac{2}{X_y^2} (4l + 1) \int_0^1 d\alpha \alpha^{2l} (1 - \alpha)^{2l} \sum_{n=0}^{\infty} \frac{-\alpha (1-\alpha)}{n! (2l+1)_n} \rho_{34}^{n} X_y^{2l+1} C_{2l}^l \left\{ \eta \bar{\eta} \right\} = \eta_{2l+1} F(2l + 1, 2l + 1; 4l + 2; \eta) - \bar{\eta}_{2l+1} F(2l + 1, 2l + 1; 4l + 2; \bar{\eta}) \quad \eta = \eta_1 - \bar{\eta}_1.
\]

(A.6)

We can sum up these expressions applying (A.3); as a result the \(\eta_1\)--dependent terms present for each \(l\) cancel and we end up with

\[
\frac{2}{\eta - \bar{\eta}} \times \sum_{l=0}^{\infty} \left\{ \eta^{2l+1} F(2l + 1, 2l + 1; 4l + 2; \eta) - \bar{\eta}^{2l+1} F(2l + 1, 2l + 1; 4l + 2; \bar{\eta}) \right\} = \frac{1}{\eta - \bar{\eta}} \left( \eta + \eta \frac{1}{\eta - \bar{\eta}} - \bar{\eta} - \frac{\bar{\eta}}{1 - \bar{\eta}} \right) = 1 + \frac{1}{(1 - \eta)(1 - \bar{\eta})} = 1 + \frac{1}{1 - \epsilon} .
\]

(A.7)

This completes the proof of Proposition 3.1. \(\Box\)

**Appendix B**  Traces of products of \(r^\mu_\nu(x)\)

We shall compute the trace of the product of tensor structures that appears in the numerator of the box diagram with propagator (3.12):

\[
B = f^{\lambda_1 \mu_1}_{\lambda_2 \mu_2} (x_{12}) f^{\lambda_2 \mu_2}_{\lambda_3 \mu_3} (x_{23}) f^{\lambda_3 \mu_3}_{\lambda_4 \mu_4} (x_{34}) f^{\lambda_4 \mu_4}_{\lambda_5 \mu_5} (x_{14}) ,
\]

\[
f^{\lambda \mu}_{\lambda' \mu'} (x) = r^\lambda_\nu (x) r^\mu_\nu (x) - r^\lambda_\mu (x) r^\mu_\nu (x) ,
\]

(B.1)

establishing on the way some useful properties of products of \(r^\mu_\nu(x) = \delta^\mu_\nu - 2 \frac{x^\mu x^\nu}{x^2 + i0} (3.4)\) (of different arguments) which appear in correlation functions of tensor fields.
We shall use repeatedly the triple product formula of [24]:

\[ r (x_{12}) r (x_{23}) r (x_{13}) = r (X_{23}), \quad \text{i.e.} \]

\[ r_\lambda (x_{12}) r_\sigma (x_{23}) r_\mu (x_{13}) = r_\lambda (X_{23}), \quad X_{23} = \frac{x_{13}}{\rho_{13}} - \frac{x_{12}}{\rho_{12}}. \quad (B.2) \]

Using the identity \( r (x)^2 = 1 \) we find

\[ R (x_{12}, x_{23}, x_{34}, x_{14}) := r (x_{12}) r (x_{23}) r (x_{34}) r (x_{14}) = \]

\[ = \left[ r (x_{12}) r (x_{23}) r (x_{13}) \right] \left[ r (x_{13}) r (x_{34}) r (x_{14}) \right] = \]

\[ = r (X_{23}) r (X_{34}). \quad (B.3) \]

where \( X_{34} \equiv X_{34}^4 = \frac{x_{14}}{\rho_{14}} - \frac{x_{13}}{\rho_{13}} \) (cf. (3.4)). Using further the relation

\[ \text{tr} \left( r (x) \right) \left( r (y) \right) = 4 \frac{(x \cdot y)^2}{x^2 y^2} + D - 4 \]

for \( D = 4 \) \quad (B.4)

we deduce (for \( \eta_i \) given by (1.5))

\[ \text{tr} \left( R (x_{12}, x_{23}, x_{34}, x_{14}) \right) = \frac{(2X_{23} \cdot X_{34})^2}{X_{23}^2 X_{34}^2} = \frac{(1 - \eta_1 - \eta_2)^2}{\eta_1 \eta_2}. \quad (B.5) \]

A simple algebra allows to reduce \( B \) (B.1) to the difference

\[ B = 8 \left\{ \left[ \text{tr} R (x_{12}, x_{23}, x_{34}, x_{14}) \right]^2 - \text{tr} \left( \left[ R (x_{12}, x_{23}, x_{34}, x_{14}) \right]^2 \right) \right\}. \quad (B.6) \]

The second term is computed using once more (B.3):

\[ \text{tr} \left( \left[ R (x_{12}, x_{23}, x_{34}, x_{14}) \right]^2 \right) = \text{tr} \left( r (X_{23}) r (X_{34}) r (X_{23}) r (X_{34}) \right) = \]

\[ = \text{tr} \left\{ r \left( \frac{\rho_{12} x_{14} - \rho_{14} x_{12}}{\rho_{24}^2} + \frac{\rho_{12} x_{13} - \rho_{13} x_{12}}{\rho_{23}^2} \right) \right. \times \]

\[ \left. \times r \left( \frac{\rho_{13} x_{14} - \rho_{14} x_{13}}{\rho_{34}^2} + \frac{\rho_{12} x_{14} - \rho_{14} x_{12}}{\rho_{24}^2} \right) \right\} = \]

\[ = \frac{(1 - 2 \eta_1 - 2 \eta_2 + \eta_1^2 + \eta_2^2)^2}{\eta_1^2 \eta_2^2}. \quad (B.7) \]

Inserting finally (B.5) and (B.7) we find

\[ B = \frac{32}{\eta_1 \eta_2} \left( 1 - 2 \eta_1 - 2 \eta_2 + \eta_1^2 + \eta_2^2 + \eta_1 \eta_2 \right). \quad (B.8) \]

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References

[1] G. Aratyunov, B. Eden, A.C. Petkou, E. Sokatchev, Exceptional non-renormalization properties and OPE analysis of chiral 4–point functions in $N = 4$ SYM, hep-th/0103230

[2] V. Bargmann, I.T. Todorov, Spaces of analytic functions on the complex cone as carriers for the symmetric tensor representations of $SO(n)$, J. Math. Phys. 18 (1977) 1141–1148

[3] K. Baumann, Bounded Bose fields in 1 + 1 dimensions commuting for space and time like distances, J. Math. Phys. 40 (1999) 1719–1737

[4] M. Bianchi, S. Kovacs, G. Rossi, Ya.S. Stanev, Properties of Konishi multiplet in $N=4$ SYM theory, hep-th/0104016

[5] H.J. Borchers, ¨Uber die Mannigfaltigkeit der interpolierenden Felder zu einer kausalen $S$–matrix, Nuovo Cim. 15 (1960) 784

[6] J. Cardy, Anisotropic corrections to correlation functions in finite size systems, Nucl. Phys. B290 (1987) 355–362

[7] N.S. Craigie, V.K. Dobrev, I.T. Todorov, Conformal invariant composite operators in QCD, Ann. Phys. (NY) 159 (1985) 411–444

[8] P. Di Francesco, P. Mathieu, D. Senechal, Conformal Field Theories, Springer, Berlin et al. 1996

[9] V.K. Dobrev, V.B. Petkova, S.G. Petrova, I.T. Todorov, Dynamical derivation of vacuum operator product expansion in Euclidean conformal quantum field theory, Phys. Rev. D13 (1976) 887–912

[10] F.A. Dolan, H. Osborn, Conformal four point functions and operator product expansion, Nucl. Phys. B599 (2001) 459–496, hep-th/0011040

[11] S. Ferrara, R. Gatto, A. Grillo, G. Parisi, Canonical scaling and conformal invariance, Phys. Lett. 38B (1972) 333; The shadow operator formalism for conformal algebra vacuum expectation values and operator products, Nuovo Cim. Lett. 4 (1972)115; General consequences of conformal algebra, in: Scale and Conformal Symmetry in Hadron Physics (SCSHOP), ed. by R. Gatto, Wiley, N.Y. 1973, pp.59–108

[12] E.S. Fradkin, M.Ya. Palchik, New developments in $d$–dimensional conformal quantum field theory, Phys. Rep. 300 (1998) 1–112; Conformal Quantum Field Theory in $D$–dimensions, Kluwer Acad. Publ. Dodrecht et al. 1996

[13] M. Grott, K.-H Rehren, On a class of bounded quantum Bose fields, Lett. Math. Phys 53 (2000) 167–179

[14] V.G. Kac, A. Radul, Quasi–free highest weight modules over the Lie algebras of differential operators on the circle, Commun. Math. Phys. 157 (1993) 125–264; Representation theory of the vertex algebra $W_{1+\infty}$, Transformation Groups 1 (1996) 41–70

[15] K. Lang, W. Ruhl, The critical $O(n)$ sigma-model at dimensions $2 < d < 4$: a list of quasi-primary fields, Nucl. Phys. B402 (12993) 573–603

[16] G. Mack, Conformal invariant quantum field theory, in SCSHP (cited in [1]) pp. 109–130; Group theoretical approach in conformal invariant quantum field theory, in: Renormalization and Invariance in Quantum Field Theory, ed Caianiello, Plenum Press, NY 1974, pp.123–157
[17] G. Mack, Convergence of operator product expansions on the vacuum in conformal invariant quantum field theory, Commun. Math. Phys. 53 (1977) 155–184

[18] G. Mack, All unitary representations of the conformal group SU(2,2) with positive energy, Commun. Math. Phys. 55 (1977) 1–28

[19] G. Mack, K. Symanzik, Currents, stress tensor and generalized unitarity in conformal quantum field theory, Commun. Math. Phys. 27 (1972) 247–281

[20] G. Mack, I.T. Todorov, Conformal invariant Green functions without ultra-violet divergences, Phys. Rev. D8 (1973) 1764–1787

[21] F. Mansouri, Dual models with global SU(2,2) symmetry, Phys. Rev. D8 (1973) 1159

[22] A.A. Migdal, Conformal invariance and bootstrap, Phys. Lett. 37B (1971) 386

[23] N.M. Nikolov, I.T. Todorov, Rationality of conformally invariant local correlation functions on compactified Minkowski space, Commun. Math. Phys. 218 (2001) 417–436 [hep-th/0009004]

[24] H. Osborn, A. Petkou, Implications of conformal invariance for field theories in general dimensions, Ann. Phys. (NY) 231 (1994) 311–362

[25] G. Parisi, L. Peliti, Calculation of critical indices, Lett. Nuovo Cim. 2 (1971) 627

[26] A.M. Polyakov, Conformal symmetry of crucial fluctuations, Zh. Eksp. Teor. Fiz., Pis. Red. 12 (1970) 538 (transl.: Sov. Phys.-JETP Lett. 12 (1970) 381)

[27] A.M. Polyakov, Nonhamiltonian approach in the conformally invariant quantum field theory, Zh. Eksp. Teor. Fiz. 66 (1974) 23–42 (transl.: Sov. Phys.-JETP 39 (1974) 10–18)

[28] K.-H. Rehren, Comments on a recent solution to Wightman axioms, Commun. Math. Phys. 178 (1996) 453–465

[29] E. Schreier, Conformal symmetry and three-point functions, Phys. Rev. D3 (1971) 981

[30] B. Schroer, A necessary and sufficient condition of the softness of the trace of the energy momentum tensor, Nuovo Cim. Lett. 2 (1971) 867; Bjorken scaling and scale invariant quantum field theories, in: SCSHS (ref. [1]] pm.43–57

[31] B. Schroer, J.A. Swieca, A.H. Volkel, Global operator product expansions in conformal invariant relativistic quantum field theory, Phys. Rev. D11 (1975) 1509–1520

[32] Ya.S. Stanev, Stress energy tensor and U(1) current operator product expansion in conformal QFT, Bulg. J. Phys. 15 (1988) 93–107

[33] R.F. Streater, A.S. Wightman, PCT, Spin and Statistics, and All That, W.A. Benmjamin, 1964; Princeton Univ. Press, Princeton, N.J. 2000

[34] K. Symanzik, On calculations in conformal invariant field theories, Nuovo Cim. Lett. 3 (1972) 734–739

[35] I.T. Todorov, Conformal invariant quantum field theory, in: W. Ruhl, A. Vancura (eds.) Strong Interaction Physics, Lecture Notes in Physics 17, Springer, Berlin et al. 1972, pp.270–299

[36] I.T. Todorov, Infinite dimensional Lie algebras in conformal QFT models, in: A.O. Barut, H.-D. Doebner (eds.) Conformal Groups and Related Symmetries. Physical Results and Mathematical Background, Lecture Notes in Physics 261, Springer, Berlin et al. 1986, pp.387–443
[37] I.T. Todorov, M.C. Mintchev, V.B. Petkova, *Conformal Invariance in Quantum Field Theory*, Scuola Normale Superiore, Pisa, 1978, 273p.

[38] A. Uhlmann, Remarks on the future tube, Acta Phys. Pol. 24, (1963) 293; The closure of Minkowski space, *ibid*. pp. 295–296; Some properties of the future tube, preprint KMU-HEP 7209 (Leipzig, 1972)