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Critical exponents of the 3d Ising and related models from conformal bootstrap

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ABSTRACT: The constraints of conformal bootstrap are applied to investigate a set of conformal field theories in various dimensions. The prescriptions can be applied to both unitary and non unitary theories allowing for the study of the spectrum of low-lying primary operators of the theory. We evaluate the lowest scaling dimensions of the local operators associated with the Yang-Lee edge singularity for $2 \leq D \leq 6$. Likewise we obtain the scaling dimensions of six scalars and four spinning operators for the 3d critical Ising model. Our findings are in agreement with existing results to a per mill precision and estimate several new exponents.

KEYWORDS: Nonperturbative Effects, Conformal and W Symmetry

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1 Introduction

Understanding conformal field theories (CFTs) is a great challenge of many branches in theoretical physics. CFTs are at the heart of critical phenomena in condensed matter physics, they control the renormalisation group flows of quantum field theories and explain the appearance of universal scaling laws [1, 2]; they even provide a tool to study quantum gravity via the AdS/CFT correspondence [3].

One of the main properties characterizing a specific CFT is the spectrum of the scaling dimensions of its local operators. With the exception of some two-dimensional CFTs (where the algebra of conformal generators is infinite-dimensional) calculating these quantities is very challenging, as they are dominated by quantum fluctuations, an effect which takes place on all length scales in these scale-invariant theories. Moreover, most CFTs are strongly coupled and difficult to study using the usual perturbative techniques of Feynman diagrams, although some of them can be accurately analysed by Monte Carlo calculations and/or strong coupling expansions [7–9].

Conformal bootstrap is a reincarnation of the bootstrap mechanism used to investigate the CFT constraints originating from the crossing symmetry of the four-point functions. It was suggested a long time ago that conformal bootstrap could give useful informations on the allowed scaling dimensions of the theory [4–6]. This idea was implemented in two-dimensional rational CFTs, i.e. those with a finite number of Virasoro primary fields. It was shown that the crossing symmetry, when combined with the modular invariance of the theory on a torus, provides the complete spectrum of scaling dimensions, modulo an integer [10].

In recent years it has been shown that the conformal bootstrap approach can give accurate predictions for specific CFTs in any space-time dimension [11–23]. The starting point of this active field of research was the observation [11] that the conformal bootstrap equations (i.e the functional equations following from the crossing symmetry) can be rewritten as an infinite system of linear homogeneous equations. In order to study the system numerically the number of unknowns and the number of equations has to be

truncated. In the simplest case, i.e. the four-point function of a single scalar field ϕ with scaling dimension Δ_ϕ , the truncated bootstrap equations take the form

$$\sum_{i=1}^N f_e^i \mathbf{p}_i = 0, \quad (e = 1, 2, \dots, M), \quad (1.1)$$

where the unknown \mathbf{p}_i is the square of the coefficient of the primary operator \mathcal{O}_i contributing to the operator product expansion (OPE) of $\phi(x)\phi(y)$. The coefficient f_e^i — made up of multiple derivatives of generalized hypergeometric functions — depends on Δ_ϕ and on the scaling dimension Δ_i of \mathcal{O}_i . Unitarity requires $\mathbf{p}_i \geq 0$; this implies that the search of solutions for eq. (1.1), when combined with a normalization condition, can be reformulated as a linear programming problem, which can be treated numerically. Its solution yields a numerical upper bound $\Delta_u = f(\Delta_\phi)$ on the dimension of the first scalar contributing to the OPE of $\phi(x)\phi(y)$ [11]. Some CFTs may exist that saturate this bound [15, 21]. Since the output of the linear programming algorithm is a solution of eq. (1.1), the low-lying spectrum of operator dimensions for these special CFTs can be evaluated. The accuracy of these calculations is particularly impressive in the case of the two-dimensional critical Ising model, where a comparison with exact results can be made. A drawback of this approach is that it can be applied only to theories saturating the unitarity bound. In addition, a physically convincing reason why some CFTs should saturate the bound is still missing.¹

Recently one of us has proposed a different approach to solve numerically the bootstrap constraints (1.1) which can be applied to a larger class of CFTs [20]. The starting point was the observation that any finite truncation of the bootstrap equations results in a set constraints on the spectrum of operator dimensions contributing to eq. (1.1), provided that the number M of equations is equal or larger than the number N of unknowns. In this case the homogeneous system admits a non-identically vanishing solution if and only if all the minors of order N are vanishing. Hence for any subset of N equations the constraint can be expressed as

$$\det \mathbf{f}_n \equiv f_n(\Delta_\phi, \Delta_1, \Delta_2, \dots, \Delta_N) = 0, \quad \left(n = 1, 2, \dots, \binom{M}{N} \right) \quad (1.2)$$

where \mathbf{f}_n is the $N \times N$ matrix of coefficients of the n th subset. As soon as M exceeds N , the number of constraints is equal or larger than N ; the all set of constraints can be used to extract the scaling dimensions of the primary operators involved, if a solution exists. If the system of constraints is over-determined, i.e. there are more independent equations than unknowns, it can be split in consistent subsystems to obtain a set of solutions; their spread provides a rough estimate of the error. A caveat: these solutions are intrinsically affected by an unknown systematic error due to the truncation in the number N of operators. Such systematic error is bound to decrease when the number of included operators increases.

This method is quite general and it can be applied to any CFT, whether unitary or not, provided the theory has a discrete spectrum and it admits a consistent truncation (1.1) of the bootstrap equations. This statement hides a subtlety that deserves to be mentioned.

¹In section 2 the condition for an actual CFT to saturate this unitarity bound is rewritten in a closed form.

The number M of homogeneous equations is not an independent parameter: indeed the larger is the number N of terms in (1.1), the better is the approximation of the r.h.s. to a constant, and hence the larger is the number M of approximately vanishing derivatives.

The method outlined above is valid only if it exists an N_0 such that $M > N_0$. If this happens for small values of N_0 we say that the CFT is easily truncable. Typical examples are the free scalar massless theories in D dimensions, where the method yields very accurate results already for $N_0 = 3$ [20]. In spite of the fact that the conformal block expansion shows good convergence properties [16], there are cases in which the assumption $M > N$ can not be consistently made for any value of N (an explicit example is shown in the next section). In these cases the associated CFT is not truncable and no numerical method based on the analysis of just one four point function can be applied.

Of course the method is predictive only if the exact form of the four-point function for a given CFT is not known. Should this be the case, we do not know a priori whether the theory is easily truncable. The symmetry properties of the model can be used to guess the structure of the quantum numbers of the low-lying primary operators. If a valid set of zeros for N determinants of the type (1.2) is found, we can say a posteriori that the CFT is easily truncable. The larger is N the smaller is the error of the estimate of the operator dimensions. Thus, once a solution is found, it can be improved by adding new operators in the game. In this way the low-lying spectrum of the Yang-Lee model in several dimensions as been evaluated in [20] and will be further improved in this paper.

A major result presented in this paper is an exact solution of the conformal bootstrap constraints (1.2) with $N = 7$ primary operators and $M = 8$ equations. It represents a consistent truncation of the four-point function $\langle \sigma(x_1)\sigma(x_2)\sigma(x_3)\sigma(x_4) \rangle$ of the 3d Ising model at the critical point described by the spontaneous breakdown of the \mathbb{Z}_2 symmetry. $\sigma(x)$ is the \mathbb{Z}_2 -odd scalar, i.e. the order parameter of the theory. The set of primary operators comprises: four \mathbb{Z}_2 -even scalars, a conserved spin 2 operator, a quasi-conserved spin 4 operator and a spin 6 operator. One of the four \mathbb{Z}_2 -even scalars is identified with the energy ε , the only relevant operator beside σ , while the other three are recurrences, $\varepsilon', \varepsilon'', \varepsilon'''$, whose dimensions are related to the critical exponents measuring the corrections to scaling. The spin 2 operator is identified with the stress tensor, while the spin 4 operator is related to the critical exponent associated with the rotational symmetry breaking. It turns out that the solution depends on a free parameter. Using the scaling dimension of the spin 4 operator as an input all the other operator dimensions can be evaluated and they differ from the best estimates for no more than 1 ‰.

Assuming this truncation is not only a solution of crossing symmetry constraints, but a part of a full-fledged CFT, we exploit some consistency checks to further enlarge the number of estimated scaling dimensions of primary operators belonging to the spectrum of the 3d critical Ising model, as reported in tables 4 and 5.

Even though a formal proof of conformality of the critical point of the 3d Ising model is still missing [24], numerical simulations and theoretical arguments leave little doubt on its validity. The conformal bootstrap approach [15] strongly supports this hypothesis. This paper, whose results are based solely on the assumption of conformal invariance of the 3d Ising critical point, further corroborates the hypothesis.

The plan of the paper is the following. In the next section we fix the notations and discuss the derivation of the homogeneous system (1.1) in relation with the notion of easily truncable CFT. Section 3 is dedicated to Yang-Lee models in various dimensions while the last section is devoted to the 3d critical Ising model and to some conclusions.

2 Conformal bootstrap constraints

The four-point function of a scalar $\phi(x)$ of a D -dimensional CFT can be parametrised as [1]

$$\langle \phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4) \rangle = \frac{g(u, v)}{|x_{12}|^{2\Delta_\phi} |x_{34}|^{2\Delta_\phi}}, \quad (2.1)$$

where Δ_ϕ is the scaling dimension of ϕ , x_{ij}^2 is the square of the distance between x_i and x_j , $g(u, v)$ is a function of the cross-ratios $u = \frac{x_{12}^2 x_{34}^2}{x_{13}^2 x_{24}^2}$ and $v = \frac{x_{14}^2 x_{23}^2}{x_{13}^2 x_{24}^2}$. $g(u, v)$ can be expanded in terms of the conformal blocks $G_{\Delta, L}(u, v)$, i.e. the eigenfunctions of the Casimir operator of $SO(D + 1, 1)$:

$$g(u, v) = 1 + \sum_{\Delta, L} \mathfrak{p}_{\Delta, L} G_{\Delta, L}(u, v), \quad (2.2)$$

with $\mathfrak{p}_{\Delta, L} = \lambda_{\phi\phi\mathcal{O}}^2$, where $\lambda_{\phi\phi\mathcal{O}}$ is the coefficient of the primary operator \mathcal{O} of scaling dimension Δ and spin L contributing to the OPE of $\phi(x)\phi(y)$.

The l.h.s. of (2.1) is invariant under any permutation of the x_i 's while the r.h.s. is not, unless $g(u, v)$ obeys the following two functional equations

$$g(u, v) = g(u/v, 1/v); \quad v^{\Delta_\phi} g(u, v) = u^{\Delta_\phi} g(v, u). \quad (2.3)$$

Since $G_{\Delta, L}(u, v) = (-1)^L G_{\Delta, L}(u/v, 1/v)$, the first equation projects out the odd spins in the expansion (2.2). The second one, after separating the identity from the other primary operators, can be rewritten as a sum rule

$$\sum_{\Delta, L} \mathfrak{p}_{\Delta, L} \frac{v^{\Delta_\phi} G_{\Delta, L}(u, v) - u^{\Delta_\phi} G_{\Delta, L}(v, u)}{u^{\Delta_\phi} - v^{\Delta_\phi}} = 1. \quad (2.4)$$

It is easy to argue that such a relation can be exactly satisfied only if the sum contains infinite terms [11]. In any finite truncation the r.h.s. of (2.4) is only approximately constant. In order to transform the sum rule into the system (1.1) of linear homogeneous equations one has to assume that the derivatives of the r.h.s. are vanishing, an assumption that in some cases cannot be made.

As an example, consider the four-point function of the energy operator : $\phi^2(x) := \lim_{\epsilon \rightarrow 0} \phi(x)^i \phi_i(x + \epsilon) - \langle \phi^j(x) \phi_j(x + \epsilon) \rangle$ in a four-dimensional $O(n)$ -invariant free field theory ($i, j = 1, 2, \dots, n$).

We have

$$g_n(u, v) = 1 + \frac{4}{n} \left(u + \frac{u}{v} + \frac{u^2}{v} \right) + \frac{u^2}{v^2} + u^2. \quad (2.5)$$

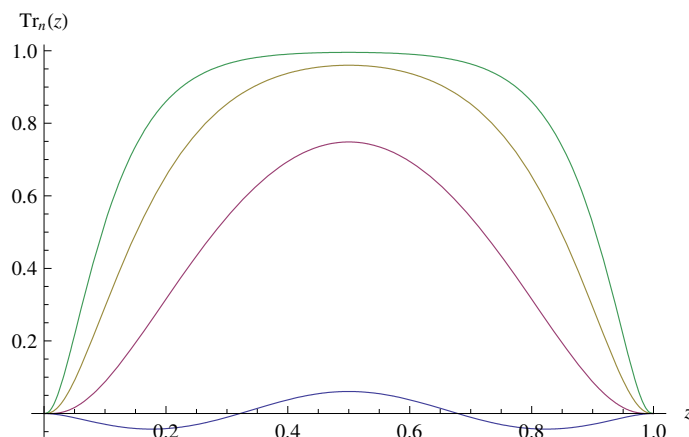


Figure 1. Progressive truncations $Tr_n(z)$ of the sum rule (2.4) at $z = \bar{z}$ for the 4-point function of the energy operator in the $4d$ $O(n)$ -invariant free massless scalar theory. The four curves represent respectively the sum of the contributions of the primaries of scaling dimensions $\Delta \leq 4$, $\Delta \leq 6$, $\Delta \leq 8$, $\Delta \leq 10$. In this specific example we have chosen $n = 1$.

The first few terms of the conformal block expansion are

$$g_n(u, v) = 1 + \frac{8}{n}G_{2,0} + \left(2 + \frac{4}{n}\right)G_{4,0} + \frac{16}{3n}G_{4,2} + \left(2 - \frac{2}{n}\right)G_{6,0} + \frac{24 + \frac{8}{n}}{5}G_{6,2} + \frac{64}{35n}G_{6,4} + \dots \quad (2.6)$$

with the parametrisation $u = z\bar{z}$ and $v = (1 - z)(1 - \bar{z})$ [25] which simplifies the functional form of the conformal blocks. In the Euclidean space z and \bar{z} are complex conjugates of each other. Choosing $z = \bar{z}$ corresponds to configurations with all four points on a circle. Figure 1 shows the first few truncations. It should be clear that these latter solutions are not flat enough around the symmetric point $z = \bar{z} = \frac{1}{2}$.

In this case it would be erroneous to guess that higher level truncations could always generate a number M of homogeneous equations sufficient to solve the system. Actually this four-point function is not truncable. The proof is surprisingly simple: if M is larger than N , then the knowledge of the operator spectrum should determine the $p_{\Delta,L}$'s, but these depend on n (see eq. (2.6)), while the spectrum of primaries of such a free field theory does not depend on it.

The easily truncable CFTs are instead those where the first few terms suffice to give an almost constant r.h.s. in (2.4). This is the case for instance of free scalar massless theories in D dimensions. Figure 2 shows the first few truncations at $D = 3$, where

$$g(u, v) - 1 \equiv \sqrt{u} + \sqrt{\frac{u}{v}} = 2G_{1,0} + \frac{1}{4}G_{3,2} + \frac{1}{64}G_{5,4} + \frac{1}{1024}G_{7,6} + \dots \quad (2.7)$$

In these cases one can follow the procedure of [11], i.e. Taylor expand (2.4) about the symmetric point $z = \bar{z} = \frac{1}{2}$ and transform a N -terms truncation of the sum rule into a set of M linear equations in N unknowns $p_{\Delta,L}$.

To be more specific, following [15] we make the change of variables $z = (a + \sqrt{b})/2$, $\bar{z} = (a - \sqrt{b})/2$ and Taylor expand around $a = 1$ and $b = 0$. It is easy to see that this

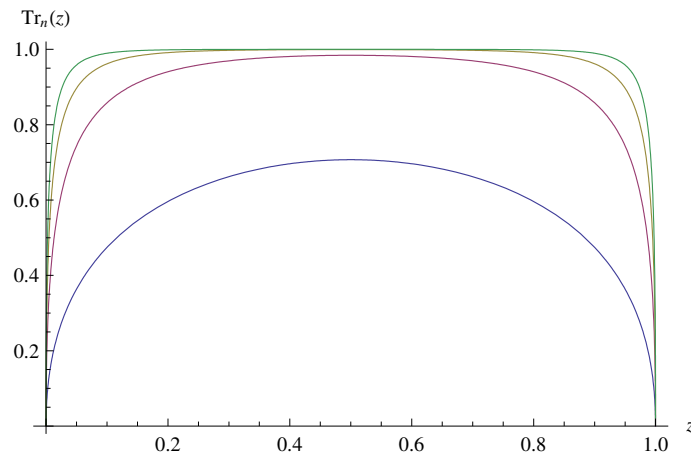


Figure 2. Progressive truncations $Tr_n(z)$ of the sum rule (2.4) for $z = \bar{z}$ of the 4-point function of a free massless scalar theory in $3d$. Few terms are sufficient to well approximate the r.h.s. of the sum rule around the symmetric point $z = \frac{1}{2}$.

expansion will contain only even powers of $(a - 1)$ and integer powers of b . The truncated sum rule can then be rewritten as one inhomogeneous equation

$$\sum_{\Delta,L} p_{\Delta,L} f_{\Delta_\phi, \Delta_L}^{(0,0)} = 1, \tag{2.8}$$

and a set of M homogeneous equations already mentioned in eq. (1.1) and here rewritten in more details

$$\sum_{\Delta,L} f_{\Delta_\phi, \Delta_L}^{(2m,n)} p_{\Delta,L} = 0, \quad (m, n \in \mathbb{N}, m + n \neq 0), \tag{2.9}$$

with

$$f_{\alpha,\beta}^{(m,n)} = \left(\partial_a^m \partial_b^n \frac{v^\alpha G_\beta(u, v) - u^\alpha G_\beta(v, u)}{u^\alpha - v^\alpha} \right)_{a=1, b=0}. \tag{2.10}$$

The number M of homogeneous equations depends on the degree of flatness of the truncation. When $M > N$ the system (2.9) becomes highly predictive, since it admits a non identically vanishing solution if and only if all the minors of order N are vanishing. The common intersection of these zeros identify the spectrum values of the scaling dimensions for the N primary operators. If there are more independent minors than unknowns we get a set of scattered solutions. Their spread gives a rough estimate of the error of the approximation. Inserting these Δ 's in eq. (2.9) and assuming that the corresponding minor has rank $N - 1$ (i.e. a simple zero) we obtain a one-parameter family of $p_{\Delta,L}$'s. The inhomogeneous equation (2.8) sets their normalization.

It is instructive to consider at the homogeneous system (2.9) from a slightly different angle. The $p_{\Delta,L}$'s can be viewed as the components of the right-eigenvector with zero eigenvalue of f . In Dirac notation we can write $f|p\rangle = 0$. Of course there is also a left-eigenvector associated with the same eigenvalue, i.e. $\langle\alpha|f = 0$. It can be noticed that $\langle\alpha|$ encodes the complete information on the low-lying spectrum of scaling dimensions of the theory. Indeed $\langle\alpha|$ is by construction orthogonal to the N columns of f . Each column

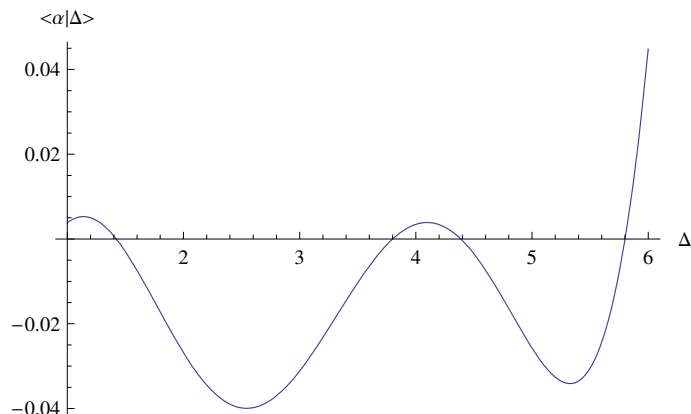


Figure 3. The scalar product $\langle \alpha | \Delta \rangle \equiv \langle \alpha | 0, \Delta, \Delta_\phi \rangle$ as a function of Δ in the critical 3d Ising model. $\langle \alpha |$ is the left eigenvector corresponding to the null eigenvalue of the 7×7 matrix f discussed in section 4 and $|\Delta \rangle$ is a column of f associated with conformal blocks of spin 0. The zeros are the scaling dimensions of the \mathbb{Z}_2 -even scalars $\varepsilon, \varepsilon', \varepsilon'', \varepsilon'''$.

depends only on a single conformal block, so it can be viewed as a vector that we denote as $|L, \Delta, \Delta_\phi \rangle$. By the knowledge of $\langle \alpha |$ it is possible to reconstruct the whole spectrum of the N primary operators as follows. Since the stress-tensor is always present among these operators, the zeros of $f(x) \equiv \langle \alpha | 2, D, x \rangle$ give the possible values of Δ_ϕ . It follows that the zeros of the scalar product $\langle \alpha | L, \Delta, \Delta_\phi \rangle$, as a function of Δ and for fixed L and Δ_ϕ , provide the whole spectrum (see an example in figure 3). In conclusion, the combination of the two eigenvectors $|\mathbf{p}\rangle$ and $\langle \alpha |$ encodes the complete information to reconstruct the (truncated) four-point function for the CFT under study.

The eigenvector $\langle \alpha |$ bears some similarity to the linear functional $\langle \Lambda |$ discussed in [11] in relation with the unitarity upper bound. In both cases the zeros of the scalar product with $|L, \Delta, \Delta_\phi \rangle$ provide the operator spectrum [14, 18], hence a solution to the crossing symmetry; on the other hand $\langle \Lambda |$ must fulfil the much stronger constraint

$$\langle \Lambda | L, \Delta, \Delta_\phi \rangle \geq 0, \quad \Delta \geq \Delta_L^o. \tag{2.11}$$

Δ_L^o is the upper unitarity bound described in the introduction: if at a given spin L the first primary operator $\mathcal{O}_{\Delta,L}$ has scaling dimension larger than this bound, the theory is not unitary. Thus if a unitary CFT saturates the bound, the scalar product above must vanish at $\Delta = \Delta_L^o$ and all the recurrences at the same L must correspond to double zeros of (2.11), i.e.

$$\langle \Lambda | L, \Delta, \Delta_\phi \rangle = 0, \quad \langle \Lambda | \frac{d}{d\Delta} | L, \Delta, \Delta_\phi \rangle = 0, \quad \Delta \in \Sigma, \Delta > \Delta_L^o, \tag{2.12}$$

where Σ denotes the spectrum of allowed operator dimensions.

It follows that $\langle \Lambda |$ can be seen as a left eigenvector of zero eigenvalue for a matrix F much larger than f , composed by the columns $|L, \Delta, \Delta_\phi \rangle$ for all $\Delta \in \Sigma$ and the columns $\frac{d}{d\Delta} |L, \Delta, \Delta_\phi \rangle$ for all recurrences. For instance, if the set of N operators of a given truncation of a CFT is composed by n_0 scalars, n_2 operators of spin 2 and so on, so that $N =$

$n_0 + n_2 + \dots + n_{2k}$, the matrix F has $2n_0 - 1$ columns for the scalars, $2n_2 - 1$ columns for the spin 2, and so on, so that the total number of columns is $N_F = 2N - k$.

Providing a solution to the crossing symmetry is equivalent to finding a N -dimensional left eigenvector $\langle \alpha | f = 0$. The additional constraint of saturating the unitarity bound requires instead to find a solution of $\langle \Lambda | F = 0$. If $\langle \Lambda |$ is unique as claimed in [18], then F is a square $N_F \times N_F$ matrix with $\det F = 0$ and $\text{rank } N_F - 1$.

In this paper we only study the zeros of $\det f$. An essential ingredient to pursuing high accuracy in these calculations is an efficient method to evaluate with high precision the conformal blocks and their derivatives. These are known in a closed form only for $D = 2, 4, 6$ [25]. The conformal blocks for generic D can be still rewritten in a closed form for $z = \bar{z}$, $L = 0$ and $L = 1$ in terms of ${}_3F_2$ hypergeometric functions

$$G_{\Delta,0}(z) = \left(\frac{z^2}{1-z} \right)^{\Delta/2} {}_3F_2 \left[\frac{\Delta}{2}, \frac{\Delta}{2}, \frac{\Delta}{2} - \alpha; \frac{\Delta+1}{2}, \Delta - \alpha; \frac{z^2}{4(z-1)} \right], \quad (2.13)$$

$$G_{\Delta,1}(z) = \left(\frac{z^2}{1-z} \right)^{\frac{\Delta+1}{2}} {}_3F_2 \left[\frac{\Delta+1}{2}, \frac{\Delta+1}{2}, \frac{\Delta+1}{2} - \alpha; \frac{\Delta}{2} + 1, \Delta - \alpha; \frac{z^2}{4(z-1)} \right], \quad (2.14)$$

with $\alpha = \frac{D}{2} - 1$.

We follow the algorithms developed in [15] which allow to write, through the use of a few recursion relations, each matrix coefficient of (2.9) as

$$f_{\Delta_\phi, \Delta_L}^{(2m, n)} = \sum_{i=1}^6 R_i(D, L, \Delta_L, \Delta_\phi) \mathcal{B}_i \left(D, \Delta_L, z = \frac{1}{2} \right), \quad (2.15)$$

where the six basis functions \mathcal{B}_i are the two conformal blocks $G_{\Delta,0}(z)$ and $G_{\Delta,1}(z)$ and their first and second derivatives with respect to z ; the R_i 's are rational functions of their arguments.

These formulas work for any value of dimension, even fractional ones. As such they have already been used to study conformal bootstrap in non-integer dimensions [22]. Here we use this property when comparing our results with the epsilon expansion of the Yang-Lee model in $6 - \epsilon$ dimensions.

3 The Yang-Lee model

The CFTs described in the present and in the next section are related to the Wilson-Fisher point of a massless scalar field model perturbed by an interacting term of the form ϕ^n , with $n = 3, 4$. In these theories the only difference between the operator content of the interacting model and that of the free field theory (the Gaussian fixed point) is in the redundant operators. Standard renormalisation group arguments are used to describe how the fusion algebra of the free theory gets modified by approaching the non-trivial fixed point (see e.g. [26]).

Besides the CFT describing its critical point, there are two other CFTs, at least, related to the 3d Ising model. One is the one-dimensional CFT associated with a line defect of this model. The low-lying spectrum of scaling dimensions of the local operators living on the

defect has been estimated using Monte Carlo simulations only recently [27]; these estimates have been supported by both epsilon expansion and conformal bootstrap calculations [23].

The other CFT is much older and originates from two seminal papers of Yang and Lee [28, 29], dated more than sixty years ago, on the analytic properties of phase transitions. It turns out that the zeros of the partition function of a ferromagnetic Ising model in D dimensions in a magnetic field h are located on the imaginary axis above a critical value $h_c(T)$ called the Yang-Lee edge singularity. Above the critical temperature $T > T_c$, in the thermodynamic limit, the density of these zeros behaves near h_c like $(h - h_c)^\sigma$, where the edge exponent σ is universal and characterizes a universality class which is not related to the spontaneous breakdown of any symmetry. The Yang-Lee universality class can be described by the non-trivial fixed point of a ϕ^3 theory [30] with imaginary coupling, thus the corresponding CFT is non-unitary hence it cannot be studied with the conformal bootstrap approach of [11].

The critical exponent σ is related to the scaling dimension of ϕ by

$$\sigma(D) = \frac{\Delta_\phi}{D - \Delta_\phi} . \tag{3.1}$$

The interest on the Yang-Lee universality class is enhanced by the discovery, in the past years, that the edge singularity is related to other quite different critical behaviours. For instance, the number per site of large isotropic branched polymers in a good solvent (i.e. *lattice animals*) obeys a power law with exponent $\varphi_I(D) = \sigma(D-2)+2$ [31], or the pressure of D -dimensional fluids with repulsive core has a singularity at negative values of activity with universal exponent $\varphi(D) = \sigma(D) + 1$ [32].

Monte Carlo simulations and other numerical methods on these systems gave accurate results for σ [33, 34]. Recent high temperature expansions of the Ising model in low magnetic field have further improved the accuracy in the whole range $2 < D < 6$ [35].

The edge exponent σ is exactly known in $D = 2$ and $D = 6$ dimensions. For the first case the exact form of the four-point function of $\phi(x)$ has been found [36]; we shall use this result to test the accuracy of our method. $D = 6$ is the upper critical dimensionality of the ϕ^3 model, above which the classical mean-field value $\sigma = \frac{1}{2}$ applies.

Our approach to conformal bootstrap only requires to know the spin and possibly other quantum numbers of the low-lying primary operators contributing to the OPE of $\phi(x)\phi(y)$. It is not necessary to know the detailed form of this OPE but simply its fusion rule that we write as

$$[\Delta_\phi] \times [\Delta_\phi] = \sum_i [\Delta_i, L_i], \tag{3.2}$$

where $[\Delta, L]$ denotes the primary operator with scaling dimension Δ and spin L and we set $[\Delta] = [\Delta, 0]$. Note that this is a shorthand way of writing the conformal block expansion (2.2) or even the first terms of a Clebsch-Gordan series of the conformal group $SO(D+1, 1)$. In a free scalar theory in D dimensions we have (see for instance (2.7))

$$[\Delta_\phi] \times [\Delta_\phi] = 1 + [\Delta_{\phi^2}] + [D, 2] + [D+2, 4] + [D+4, 6] + \dots, \tag{3.3}$$

where $[D+L-2, L]$ is a conserved spinning operator. In an interacting theory only the spin 2 — the stress-tensor $T_{\mu\nu}$ — is generally conserved. Moreover in the ϕ^3 theory

2d Yang Lee model — Truncated Fusion Rules					
	[Δ_ϕ] \times [Δ_ϕ] = 1 + [Δ_ϕ] + [2, 2] + [4] + [$\Delta_4, 4$]				Exact results
		+ [$\Delta_6, 6$] + [6, 6]	+ [$\Delta_6, 6$] + [Δ']	+ [6, 6] + [Δ']	
Δ_ϕ	-0.385(7)	-0.40033(1)	-0.40062(3)	-0.39777(2)	$-\frac{2}{5}$
Δ_4	3.70(4)	3.5904(3)	3.58961(5)	3.5963(2)	$\frac{18}{5}$
Δ_6	—	5.593(3)	5.590(3)	—	$\frac{28}{5}$
Δ'	—	—	7.60(1)	7.61(8)	$\frac{38}{5}$
c	-4.5(1)	-4.38(1)	-4.38(1)	-4.44(1)	$-\frac{22}{5}$
\mathbf{p}_ϕ	-3.67(1)	-3.6524(3)	-3.6524(3)	-3.6557(3)	-3.65312...

Table 1. The low-lying spectrum of primary operators of the 2d Yang-Lee model estimated for 4 different truncations of the fusion rule (3.4) and compared with the exact results (last column). A rough estimate of the errors is obtained by the spread of the solutions generated by different set of equations. An example is depicted in figure 4.

ϕ^2 is a redundant operator, since at the non-trivial fixed point in $D < 6$ dimensions it is proportional to $\partial^2\phi$ by the equation of motion. Thus ϕ^2 and its derivatives become descendant operators of the only relevant operator ϕ of the Yang-Lee universality class. Thus the fusion rule that characterizes the low-lying spectrum of this CFT is expected to be

$$[\Delta_\phi] \times [\Delta_\phi] = 1 + [\Delta_\phi] + [D, 2] + [\Delta_4, 4] + [\Delta'] + \dots, \tag{3.4}$$

where [Δ'] in the perturbative renormalisation group analysis corresponds to ϕ^3 . This scalar cannot be neglected if we include in the fusion rule the spin 4 operator, since in $6 - \epsilon$ dimensions $\Delta_{\phi^3} < \Delta_4$. Of course we expect also a term [$\Delta_6, 6$], like in a free-field theory, however its contribution becomes important only in the two-dimensional case. We must also add a scalar of dimension 4 corresponding to $T\bar{T}$.

Comparison with the exact solution at $D = 2$ shows that there are actually two primary operators of spin 6. One of them is conserved and coincides with T^3 . Inserting truncations of such a fusion rule in eq. (1.2) we found only isolated solutions where the number M of fulfilled equations is larger than the number of unknowns Δ 's, thus this CFT is easily truncable. To solve numerically these equations we used an iterative Newton-Raphson method, which proves to be very effective.² In all the cases considered we worked with $M = 7$ homogeneous equations. The number of unknown Δ 's vary from 2 to 4 in the $D = 2$ case (see table 1), while for $D > 2$ we have the 3 unknowns Δ_ϕ, Δ_4 and Δ' .

In table 1 we report the results of a truncation of this fusion rule to the first 4 terms and three different truncations of 6 terms. As expected the accuracy increases with the number of included conformal blocks. Inserting these Δ 's in the linear system (2.8) and (2.9) we can evaluate the OPE coefficients, in particular \mathbf{p}_{Δ_ϕ} that turns out to be very close to the known exact result. Similarly from $\mathbf{p}_{2,2}$ we can extract the central charge $c = \Delta_\phi^2/\mathbf{p}_{2,2}$.

²This method is implemented for instance in the `FindRoot` function of Mathematica.

D dimensional Yang-Lee model — the edge exponent σ					
D	bootstrap	Ising in H	Fluids	Animals	ϵ -expansion
2	-0.1664(5)	-0.1645(20)	-0.161(8)	-0.165(6)	(exact -1/6)
3	0.085(1)	0.077(2)	0.0877(25)	0.080(7)	0.079-0.091
4	0.2685(1)	0.258(5)	0.2648(15)	0.261(12)	0.262-0.266
5	0.4105(5)	0.401(9)	0.402(5)	0.40(2)	0.399-0.400
6	0.4999(1)	0.460(50)	0.465(35)	—	1/2

Table 2. Estimates of $\sigma(D)$ from the truncated fusion rule (3.4) compared with the best estimates from strong coupling expansions [35], fluids [33] and lattice animals [34]. A more complete list of estimates can be found in [35]. The bootstrap result for $D = 2$ is calculated from tab 1.

D dimensional YL model — Δ_4 and Δ'			
D	p_ϕ	Δ_4	Δ'
3	-3.88(1)	4.75(1)	5.0(1)
4	-2.72(1)	5.848(1)	6.8(1)
5	-0.95(2)	6.961(1)	6.4(1)

Table 3. Estimates of Δ_4 and Δ' and the OPE coefficient p_ϕ from the truncated sum rule (3.4). These are, to the best of our knowledge, the only known estimates for these quantities.

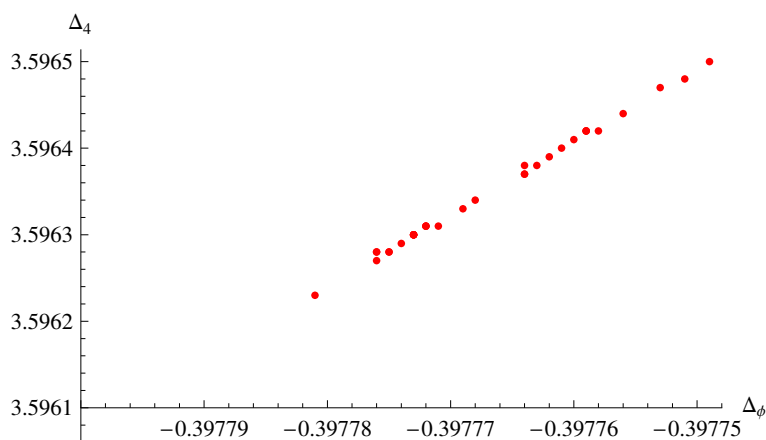


Figure 4. Dispersion of the estimates of Δ_4 versus Δ_ϕ in the 2d Yang-Lee model. Notice the strong correlation between these two quantities.

We encountered isolated solutions also in $D > 2$. The estimates for $D \leq 6$ generated by the truncated fusion rule (3.4) are reported in tables 2 and 3.

As we have already noticed, the present method works for any value of dimension. As such it is interesting to follow the flow of the spectrum from the six-dimensional free-field theory in $6 - \epsilon$ dimensions. The epsilon expansion of $\sigma(D)$ known up to, and including, ϵ^3

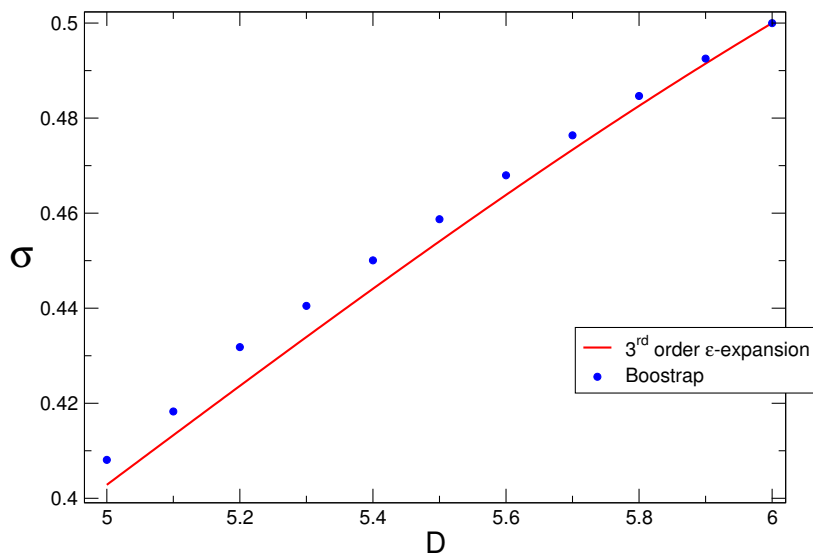


Figure 5. Plot of the function $\sigma(D)$ given by eq. (3.5) compared with the results of our bootstrap analysis for $5 \leq D \leq 6$. In this range the different resummations of the epsilon expansion do not give visually different results. The difference between the bootstrap estimates and the epsilon expansion are probably due to the contributions of higher order in ϵ .

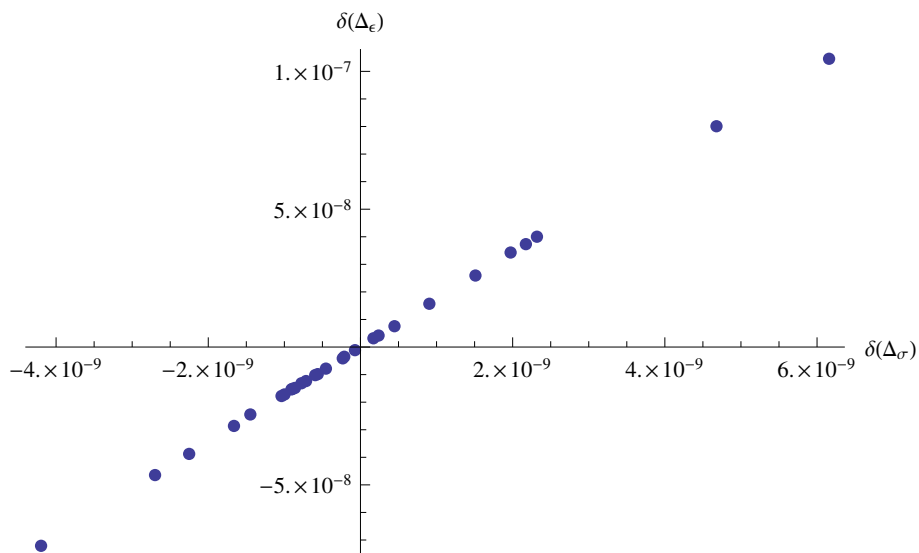


Figure 6. Dispersion of the solutions in the plane $(\Delta_\sigma, \Delta_\epsilon)$ at $\Delta_4 = 5.022$. The origin coincides with the mean value of these quantities. Notice the microscopic value of the spread.

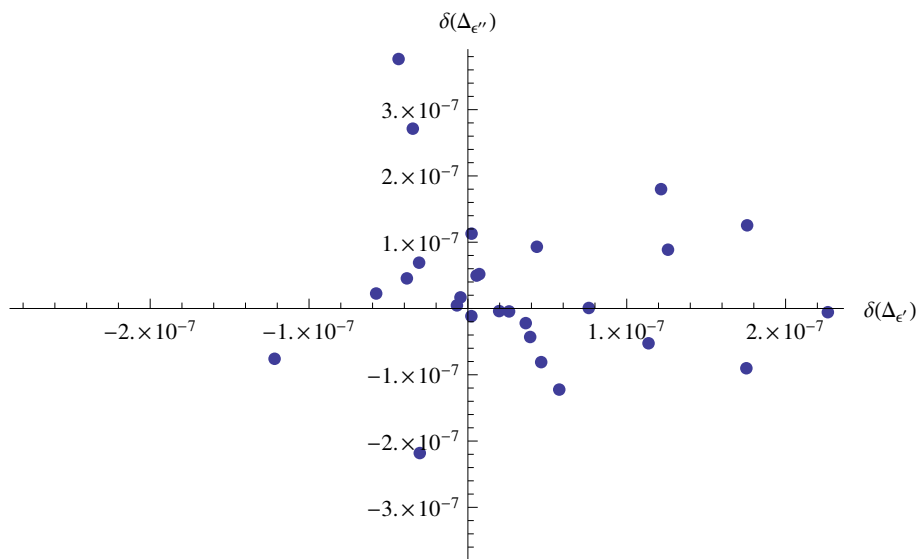


Figure 7. Dispersion of the solutions in the plane $(\Delta_{\epsilon'}, \Delta_{\epsilon''})$ at $\Delta_4 = 5.022$. The origin coincides with the mean value of these quantities.

contributions is [37]

$$\sigma(6 - \epsilon) = \frac{1}{2} - \frac{1}{12}\epsilon - \frac{79}{3888}\epsilon^2 + \left(\frac{\zeta(3)}{81} - \frac{10445}{1250712} \right) \epsilon^3 + O(\epsilon^4). \quad (3.5)$$

In figure 5 we plot this function in the range $5 \leq D \leq 6$ as well as the bootstrap estimates.

4 3d critical Ising model

We investigate the 3d critical Ising model with the same strategy applied to the Yang-Lee model, namely we start from the fusion rules of the free-field theory in the upper critical dimension — $D = 4$ for a ϕ^4 theory — and see how this is modified in the non-trivial fixed point in $4 - \epsilon$ dimensions. This time the first redundant operator is ϕ^3 which does not contribute to the OPE of $\phi(x)\phi(y)$, being an odd operator, so no further insight to simplify the fusion rule (3.3) can be obtained. If we include the spin 4 operator, we cannot neglect the leading irrelevant scalar primary ϵ' , corresponding to ϕ^4 , since $\Delta_{\phi^4} < \Delta_4$ in $4 - \epsilon$ dimensions.

A first solution of the bootstrap constraints (1.2) is associated with the truncated fusion rule

$$[\Delta_\sigma] \times [\Delta_\sigma] = 1 + [\Delta_\epsilon] + [\Delta_{\epsilon'}] + [3, 2] + [\Delta_4, 4], \quad (4.1)$$

with $\Delta_\sigma \equiv \Delta_\phi$. These scaling dimensions are related to the conventional critical exponents by $\eta = 2\Delta_\sigma - 1$, $\nu = 1/(3 - \Delta_\epsilon)$, $\omega = \Delta_{\epsilon'} - 3$ and $\omega_{NR} = \Delta_4 - 3$, where η is the magnetic exponent, ν the thermal exponent, ω is the exponent controlling the scaling corrections of rotationally invariant operators and finally ω_{NR} is the one controlling the correction to scaling of non-rotationally invariant operators (see [38], section 1.6.4).

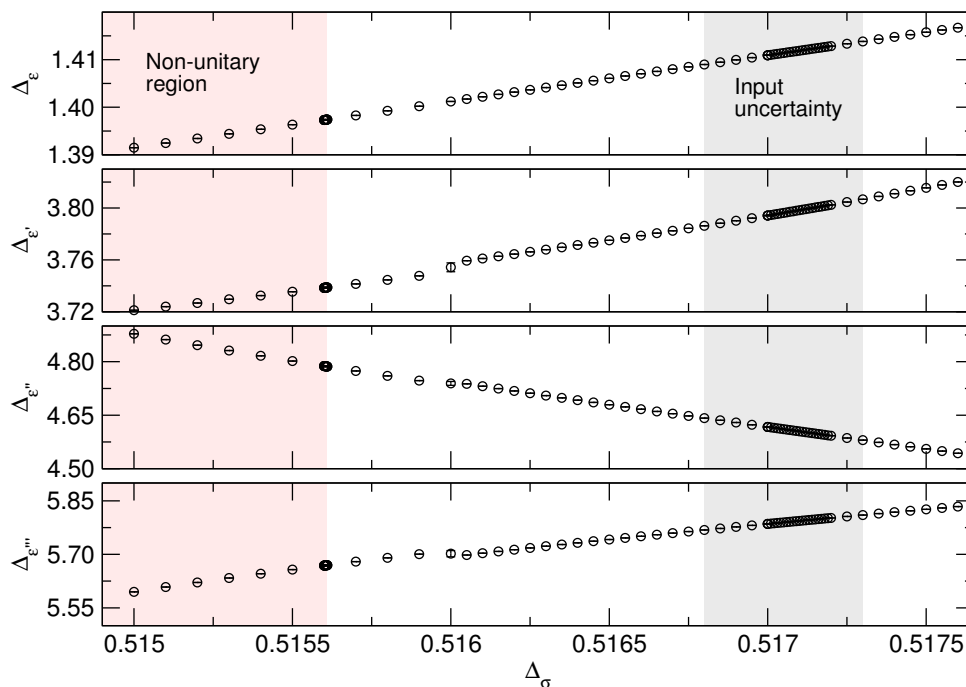


Figure 8. Spectrum of the scalar operators of the one-dimensional family of solutions for the fusion rule (4.2). The transition to the non-unitary region (the shadowed region on the left) is characterised by a change of sign of the coupling $\rho_{\epsilon''}$ (see figure 9). Presumably, increasing the number of primaries in our analysis, the gap between the two shaded regions should shrink, in accordance with the conjecture that the critical Ising model saturates the unitarity bound [15].

It turns out that the bootstrap constraints (1.2) applied to the $M = 5$ homogeneous equations (2.9) with $m + n \leq 2$ are fulfilled. In this case there is not an isolated solution like in Yang-Lee model, but a one-parameter family of solutions. This result stems from the following: the Newton-Raphson method requires an initial guess of the solution, then a sequence of approximate solutions is generated rapidly converging to an accurate one. In the present case it turns out that a small variation of the initial guess yields a small change to the final point of the sequence, showing that there is indeed a family of solutions. In order to verify that this family is one-dimensional we kept the scaling dimension Δ of one of the unknowns as a free external parameter and we verified that for any given value of Δ the solution does not depend any longer on the initial guess. By varying the input parameter Δ we generated the one-dimensional family drawn in figures 8 and 9.

Using as input $\Delta_4 \equiv \omega_{NR} + 3 = 5.0208(12)$, which is the most precisely known scaling dimension of the 3d critical Ising model (see [38], section 3.2.1), we obtain a first rough estimate of the scaling dimensions of the operators involved in (4.1), namely $\Delta_\sigma = 0.5145(3)$, $\Delta_\epsilon = 1.3735(15)$, $\Delta_{\epsilon'} = 3.80(1)$.

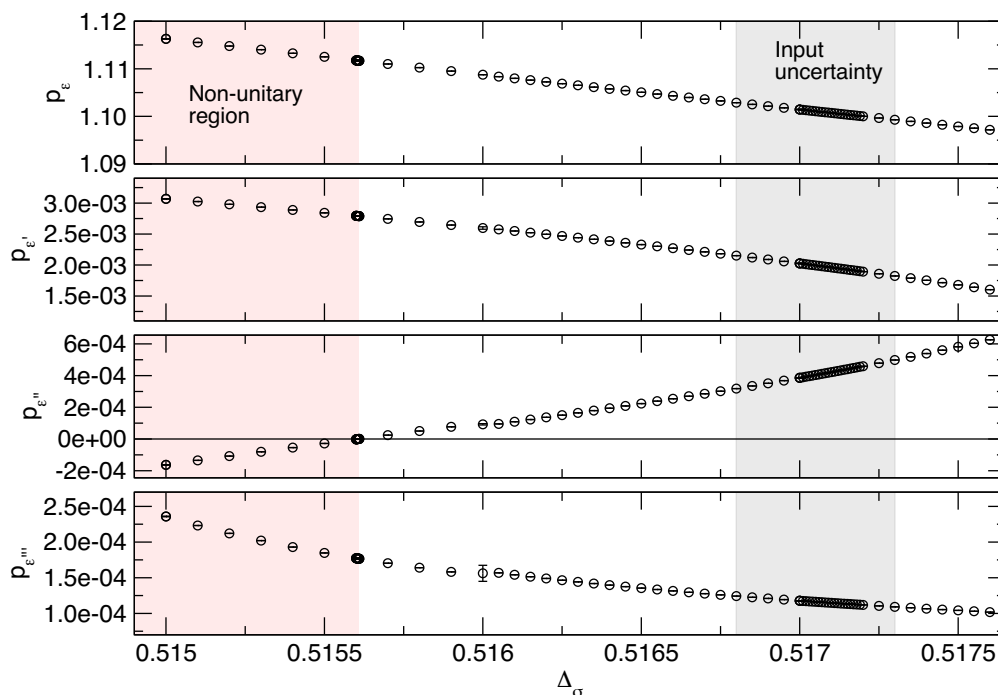


Figure 9. Evolutions of the couplings of the scalar operators to the $\sigma\sigma$ channel for the fusion rule (4.2) along the one-dimensional family of solutions. The transition to the non-unitary region (the shadowed region on the left) is characterised by a change of sign of the coupling $p_{e''}$.

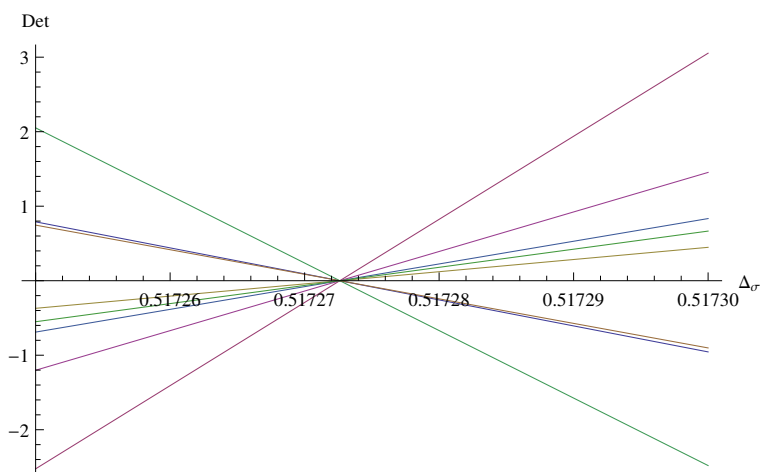


Figure 10. The zeros of the eight determinants associated to the solution (4.2) at $\Delta = \Delta_\sigma$, keeping all other scaling dimensions fixed at the values of the exact solution. In view of the microscopic spread of the solutions shown in figure 6, these different solutions are visually indistinguishable.

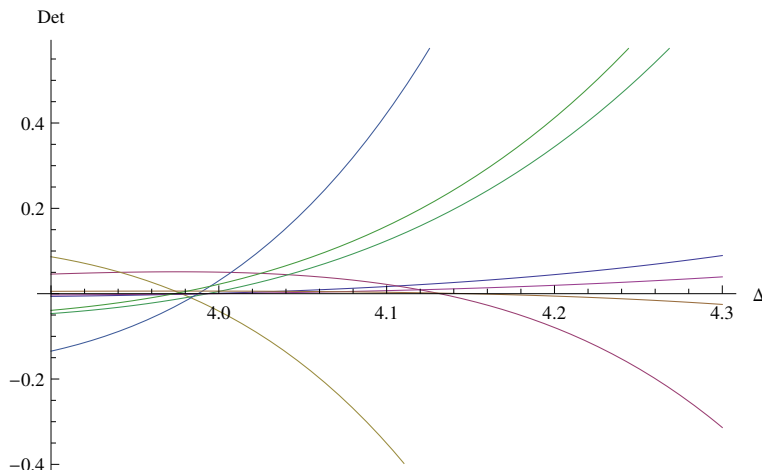


Figure 11. Plot of the same determinants of figure 10 in the region $\Delta \simeq 4$ where there is an approximate solution for Δ_σ . The spread of the zeros gives a rough estimate of the error.

A better approximation is obtained including also a spin 6 operator, and, since in the ϵ expansion the scaling dimensions of the fields ϕ^6 and ϕ^8 are smaller than Δ_6 , also these need to be included, resulting in the next to leading scalars ϵ'' , ϵ''' .

Precisely we found that the truncated fusion rule

$$[\Delta_\sigma] \times [\Delta_\sigma] = 1 + [\Delta_\epsilon] + [\Delta_{\epsilon'}] + [\Delta_{\epsilon''}] + [\Delta_{\epsilon'''}] + [3, 2] + [\Delta_4, 4] + [\Delta_6, 6], \quad (4.2)$$

admits an exact solution of the bootstrap constraints (1.2) applied to the set of $M = 8$ homogeneous equations (2.9) with $m + n \leq 3$ and $n < 3$. The unknowns are the six quantities $\Delta_\sigma, \Delta_\epsilon, \Delta_{\epsilon'}, \Delta_{\epsilon''}, \Delta_{\epsilon'''}, \Delta_6$, so we can split the set of homogeneous equations in 28 consistent subsystems and as many bootstrap constraints (1.2) given by the vanishing of determinants of 7×7 matrices. The corresponding spread of the solutions is far smaller (see figures 6 and 7) than the error of the input parameter Δ_4 . In figure 8 the one-parameter family of solutions is depicted. Below $\Delta_\sigma = 0.515607$ this line enters a non-unitary region characterised by a change of sign in $p_{\epsilon''}$. It is worth reporting that close to this family there are other one-parameter solutions which form an intricate and thick web in regions of the space of operator dimensions. Some of these lines intersect or bifurcate. Most of them correspond to non-unitary CFTs.

There is a number of consistency checks that can be exploited to enlarge the spectrum of estimated scaling dimensions. We expect in particular that the leading irrelevant odd scalar σ' fulfils the same fusion rule as σ , at least at this truncation level.³ This implies that if we treat Δ_σ as a free parameter Δ , keeping all the other scaling dimensions fixed, the 8 determinants which vanish at $\Delta = \Delta_\sigma$ should also be approximately zero at $\Delta = \Delta_{\sigma'}$, and this actually happens in figures 10 and 11.

Similarly, we expected that the fusion rule of the energy operator $\epsilon(x)$ should coincide at this truncation level with that of $\sigma(x)$. However we did not find a corresponding solution.

³Since ϕ^3 is a redundant operator, this primary operator is associated with ϕ^5 in the renormalisation group analysis in $4 - \epsilon$ dimensions.

scalar operators	σ	σ'	ε	ε'	ε''	ε'''
Δ , best estimates	0.51813(5)	$\gtrsim 4.5$	1.41275(25)	3.84(4)	4.67(11)	—
Δ , bootstrap	0.51705(25)	4.05(5)	1.4114(24)	3.796(10)	4.61(3)	5.79(2)

Table 4. Low-lying scalar primary operators and their scaling dimensions obtained by the solution of the conformal bootstrap constraints (1.2). The comparison values are taken from [9] and [38]. Recent calculations in the functional renormalisation group approach (see [39] and references therein) predict much higher values for $\Delta_{\varepsilon''}$ and $\Delta_{\varepsilon'''}$ (see table 7). However they also give $\eta = 0$, implying that in their case $\sigma(x)$ is a free field.

spinning operators	$[\Delta_2, 2]$	$[\Delta'_2, 2]$	$[\Delta_4, 4]$	$[\Delta'_4, 4]$	$[\Delta_6, 6]$
Δ , best estimates	—	—	5.0208(12)	—	7.028(8)
Δ , bootstrap	5.117(1)	6.20(1)	input	6.70(1)	7.065(3)

Table 5. Low-lying primary operators with spin and their scaling dimensions obtained by the solution of the conformal bootstrap constraints (1.2). The input value for Δ_4 is taken from [38]. The estimate for the spin 6 operator is extracted from [19].

We found instead an approximate solution for the following enlarged fusion rule

$$[\Delta_\varepsilon] \times [\Delta_\varepsilon] = 1 + [\Delta_\varepsilon] + [\Delta_{\varepsilon'}] + [\Delta_{\varepsilon''}] + [\Delta_{\varepsilon'''}] + [3, 2] + [\Delta_2, 2] + [\Delta'_2, 2] + [\Delta_4, 4] + [\Delta'_4, 4] + [\Delta_6, 6]. \tag{4.3}$$

Here, besides all the operators of (4.2), contributions appear from three new spinning operators, namely two recurrences of spin 2 and one recurrence of spin 4. In this case the conformal bootstrap constraints are associated with determinants of 10×10 matrices, the number of homogeneous equations considered is $M = 10$, while the number of unknown Δ 's is 8 (remember that Δ_4 is set as external input). In order to have a sufficient number of homogeneous equations we also considered the fusion rule (4.3) where in the l.h.s. we replaced ε with ε' . Unlike the solution of (4.2), where the accuracy of the zeros was of the order 10^{-16} , here we found only approximate zeros of the order $10^{-2} - 10^{-3}$. It is important to notice that the estimate of scaling dimensions of the operators appearing in both fusion rules almost coincide. The resulting set of estimates for the whole set of primary operators considered in our analysis is reported in tables 4 and 5. Finally, inserting these values in the linear systems (2.8) and (1.1) we can extract the corresponding OPE coefficients. They are reported in figure 9 and in table 6. A word of caution: our data are affected by an unknown systematic error due to the truncation in the number of operators. Such systematic error is bound to decrease when an higher number of operators is used. The errors reported in the tables only take into account the spread of the solutions and the uncertainty of the input parameter.

It is worth noting that, beside contributing to the to the n -point functions, the scalar operators also play an important role in the description of the thermodynamic functions at criticality. For instance the critical behaviour of the magnetic susceptibility in terms of

p_ε	$p_{3,2}$	$p_{\varepsilon'}$	$p_{\varepsilon''}$	$p_{\varepsilon'''}$	p_{Δ_4}	p_{Δ_6}
1.101(2)	0.2855(4)	0.0019(2)	0.00041(1)	0.000116(7)	0.01612(2)	0.001541(2)
1.1117	0.28326	0.0027906	0.4(1.3)E-8	0.0001768	0.01601	0.0015300

Table 6. In the first row we report the values of the couplings related to the results presented in tables 4 and 5 for the fusion rule (4.2). The error originates from the uncertainty on the input parameter. In the second row we report the values of the couplings for the point $\Delta_\sigma = 0.515607$, which lies at the boundary of the unitarity region (see figure 8). The error due to the spread of the solutions is negligible (it ranges from 10^{-13} to 10^{-8}) hence it is not reported, except for $p_{\varepsilon''}$. The central charge extracted from the first row is $c/c_{free} = 0.9364(22)$; while the second row we obtain $c/c_{free} = 0.93853(1)$. Presumably, increasing the number of included primaries in our analysis, the gap between the point at the boundary and the Ising point should shrink, in accordance with the conjecture that the critical Ising model saturates the unitarity bound [15].

method	η	ν	ω	ω_2	ω_3
this work	0.0341(5)	0.629(1)	0.80(1)	1.61(3)	2.79(2)
HT [8]	0.03639(15)	0.63012(16)	0.83(5)	—	—
MC [9]	0.03627(10)	0.63002(10)	0.832(6)	—	—
RG [39]	—	0.3–0.69	0.60–0.76	2.8–3.7	5.2–7.0
SF [40]	0.040(7)	0.626(9)	0.855(70)	1.67(11)	—
UB [41]	0.03631(3)	0.62999(5)	0.8303(18)	4	7.5

Table 7. Leading exponents and leading and subleading correction-to-scaling exponents obtained using high temperature (HT) expansions, Monte Carlo (MC) calculations, functional renormalization group (RG), scaling field (SF) methods and unitarity bounds (UB) on solutions of crossing symmetry compared with the results of this work.

the reduced temperature t is given by the Wegner expansion

$$\chi = t^{-\gamma}(a_{0,0} + a_{0,1}t + a_{0,2}t^2 + \dots + a_{1,1}t^{\omega\nu} + a_{1,2}t^{2\omega\nu} + \dots + a_{2,1}t^{\omega_2\nu} + \dots + a_{3,1}t^{\omega_3\nu} + \dots),$$

where the leading exponent γ is not directly evaluated in the bootstrap approach, but can be expressed in terms of ν and η through $\gamma = \nu(2 - \eta)$. The leading and subleading correction-to-scaling exponents are related to the dimensions of the scalar operators by $\Delta_{\varepsilon'} = 3 + \omega$, $\Delta_{\varepsilon''} = 3 + \omega_2$ and $\Delta_{\varepsilon'''} = 3 + \omega_3$. The determinations of these exponents in our work are compared with other determinations in table 7.

In conclusion, in this paper we have developed a general method implementing the crossing symmetry constraints in a large class of CFT, both unitary and not. Starting with the knowledge of the fusion rules, this method can generate systematically the correlation functions, by searching for the zeros of certain $N \times N$ determinants made up of multiple derivatives of generalized hypergeometric functions. Such a study can be performed with the use of modest computing power, indeed all the results here presented have been obtained by using a single workstation. The method’s application to the Yang -Lee edge singularity

for $2 \leq D \leq 6$ and to the 3d critical Ising model gives rather accurate results and can be further improved by enlarging the number N of primary operators included in the analysis.

Added note. After this work was completed, we became aware of the preprint [41] where the low-lying spectrum of the 3d critical Ising model is precisely calculated under the assumption that this model has the smallest central charge, by analogy with the two-dimensional case. In this way very accurate values of the scaling dimensions Δ_σ and Δ_ε are obtained, which are more precise than our estimates with the determinant method.

In the quoted paper strong numerical evidence is reported, showing that some operators disappear from the spectrum as one approaches the putative 3d Ising point. Our method, which is able to follow our one-parameter family of solutions also in the non-unitary region (see figure 8), provides a simple explanation of this phenomenon. Moving across the two regions at least one of our couplings changes sign. It follows that the corresponding operator decouples at the boundary of the unitarity region, where the putative critical Ising point should be. In our solution we can see the vanishing of the coupling of one operator (see table 6) and presumably, enlarging the number of primaries included in the analysis, the number of decoupled operators at the boundary should increase. The operator that decouples is the scalar at $\Delta \simeq 4.60$ that we identify with ε'' . Despite this decoupling, this operator does not disappear from the spectrum of the CFT at the boundary since its coupling to the $\varepsilon\varepsilon$ channel is still different from zero. This operator is not present in the analysis of [41]. A possible explanation for such a discrepancy could be that the coupling of this operator is rather small along the entire one-parameter family of solutions (see fig 9), and the method of [41] cannot detect operators with too small couplings.

Acknowledgments

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