Kac and New Determinants for
Fractional Superconformal Algebras

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Abstract

We derive the Kac and new determinant formulae for an arbitrary (integer) level $K$ fractional superconformal algebra using the BRST cohomology techniques developed in conformal field theory. In particular, we reproduce the Kac determinants for the Virasoro ($K = 1$) and superconformal ($K = 2$) algebras. For $K \geq 3$ there always exist modules where the Kac determinant factorizes into a product of more fundamental new determinants. Using our results for general $K$, we sketch the non-unitarity proof for the $SU(2)$ minimal series; as expected, the only unitary models are those already known from the coset construction. We apply the Kac determinant formulae for the spin-4/3 parafermion current algebra (i.e., the $K = 4$ fractional superconformal algebra) to the recently constructed three-dimensional flat Minkowski space-time representation of the spin-4/3 fractional superstring. We prove the no-ghost theorem for the space-time bosonic sector of this theory; that is, its physical spectrum is free of negative-norm states.

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I. INTRODUCTION

The Kac determinant [1] for the Virasoro algebra is a very useful tool for analyzing the minimal unitary series in CFT [2] and for understanding the no-ghost theorem in the bosonic string theory [3]. So far, the Kac determinants have been known only for the Virasoro and superconformal algebras [4]. These algebras are special cases of the so-called fractional superconformal algebras (FSCAs) [5,6] labeled by the level $K$ of the $SU(2)_K$ current algebra [7] (The level $K = 1$ and 2 FSCAs are precisely the Virasoro and superconformal algebras). In this paper we derive the Kac and new determinant formulae for an arbitrary (integer) level $K$ FSCA.

In section II we classify the modules of the level $K$ FSCA and present the Kac determinant formulae for each module. For $K \geq 3$ there always exist modules where the Kac determinant factorizes into a product of more fundamental new determinants. For these modules we also present the new determinant formulae. We recover the well-known Kac determinant formulae for the Virasoro and superconformal algebras, and elaborate on the case of the spin-$4/3$ parafermion current algebra (i.e., the $K = 4$ FSCA) of Zamolodchikov and Fateev [8]. This algebra has three modules (the so-called $S$-, $D$- and $R$-modules) as opposed to the superconformal algebra that has only two modules (i.e., the Neveu-Schwarz and Ramond modules). We discuss in detail the Kac determinant formulae for the $S$- and $D$-modules and the Kac and new determinant formulae for the $R$-module. We also give a few low lying levels to clarify notations.

In section III we derive the Kac and new determinant formulae using a generalization [9,10] of the BRST operator of Felder [11]. We also derive the relation between the Kac and new determinants. The zeros of the Kac determinants were given in Ref [5,9,10]. To derive the Kac and new determinant formulae, we need to fix the orders of these zeros, as well as the zero mode contribution. For level $K \geq 3$ FSCAs, the orders of some zeros of the Kac and new determinants for a given module come from the string functions that give the counting of states in a different module. This is never the case for the Virasoro and superconformal algebras.

Next, we turn to the applications of the Kac determinant formulae. In section IV, using our results for general $K$, we deduce the values of the central charge $c$ and conformal dimension $h$ for which FSCAs can have unitary representations. In particular, we sketch the non-unitarity proof for the $SU(2)$ minimal series and argue that the only unitary models are precisely the known $SU(2)_K \otimes SU(2)_L/SU(2)_{K+L}$ coset models [12].

Recent evidence supports the existence of the so-called fractional superstrings (FSS) [13]. A particular representation (with $c = 5$) of the spin-$4/3$ FSS has bosons and fermions living in three-dimensional flat Minkowski space-time [14]. In section V we prove the no-ghost theorem for the space-time bosonic sector of this theory; using the Kac determinant formulae for the $K = 4$ FSCA, we generalize Brower and Thorn’s old no-ghost theorem [3] and argue that the physical state conditions remove all negative-norm states.

We also use the Kac determinant formulae for the $K = 4$ FSCA to examine the physical null state structure of the critical spin-$4/3$ FSS ($c = 10$) and find extra sets of zero-norm physical states; this property is expected of any consistent string theory.
II. KAC AND NEW DETERMINANTS

In this section we present the Kac and new determinant formulae for an arbitrary level \( K \) fractional superconformal algebra (FSCA). We recover the well-known Kac determinants for the Virasoro \((K = 1)\) and superconformal \((K = 2)\) algebras, and elaborate on the case of the spin-4/3 parafermion current algebra \(\text{i.e., the } K = 4 \text{ FSCA}\) of Zamolodchikov and Fateev \[8\]. We derive the Kac and new determinant formulae in the next section.

A. Preliminaries

By definition the Kac determinants for a given algebra are representation independent. The simplest representation of the level \( K \) FSCA (convenient for derivation of the Kac and new determinants) is a non-interacting system of the \( \mathbb{Z}_K \) parafermion (PF) \[15\] and a free boson. The \( SU(2)_K/U(1) \) coset model is a realization of the \( \mathbb{Z}_K \) PF theory. The chiral \( SU(2)_K \) WZW theory \[7\] has central charge

\[
c_0 = \frac{3K}{K + 2}
\]

and consists of holomorphic Virasoro primary fields \( \Phi^j_m(z) \) of conformal dimensions \( \frac{j(j+1)}{K+2} \). The indices \( j, m \in \mathbb{Z}/2 \) label \( SU(2) \) representations, where

\[
0 \leq j \leq K/2 \text{ and } |m| \leq j \text{ with } j - m \in \mathbb{Z} .
\]

When we factor a \( U(1) \) subgroup out of \( SU(2)_K \), we correspondingly factor the primary fields as

\[
\Phi^j_m(z) = \phi^j_m(z) \exp \left\{ i \frac{m}{\sqrt{K}} \varphi(z) \right\} .
\]

Here \( \varphi \) is the free \( U(1) \) boson normalized so that \( \langle \varphi(z)\varphi(w) \rangle = -2 \ln(z-w) \). \( \phi^j_m(z) \) are the Virasoro primary fields in the \( \mathbb{Z}_K \) PF theory with conformal dimensions

\[
\Delta^j_m = \frac{j(j+1)}{K+2} - \frac{m^2}{K} \text{ for } |m| \leq j .
\]

The \( SU(2)_K \) algebra has an automorphism

\[
\phi^j_m = \phi^{K/2-j}_{m+K/2} .
\]

With these identifications, we can consistently extend \( \phi^j_m \) to \( |m| > j \).

The fusion rules of the PF fields follow from those of the \( SU(2)_K \) theory:

\[
[\phi^j_{m_1}] \otimes [\phi^j_{m_2}] \sim \sum_{j=|j_1-j_2|}^r [\phi^j_{m_1+m_2}] ,
\]
where \( r = \min\{j_1 + j_2, K - j_1 - j_2\} \).

The \( SU(2)_K \) currents are given by
\[
J^\pm = \sqrt{K} \phi^0_{\pm 1} \cdot \exp(\pm i \varphi / \sqrt{K}) ,
\]
\[
J^0 = \frac{i}{2} \sqrt{K} \partial \varphi ,
\]
where \( \phi^0_{\pm 1} \) have conformal dimension \((K - 1)/K\).

The \( SU(2)_K \) primary fields read
\[
\Phi^j_j(z) = \varphi^j_j(z) \cdot \exp(i j \varphi(z) / \sqrt{K}) .
\]

Note \( \Phi^j_{-j} = (\Phi^j_j)^\dagger \) are also primary.

Next we consider the so-called fractional supercurrent (whose conformal dimension is \( \Delta = (K + 4)/(K + 2) \))
\[
G = J^a_{-1} \Phi^1_a = \frac{\sqrt{K}}{2} \epsilon \partial \varphi + \frac{iK}{K + 4} \eta .
\]

Here \( J^a_{-1} \) are the conformal dimension 1 creation modes of the \( J^a(z) \) currents; \( \epsilon \equiv \phi^0_1 \) is the so-called PF energy operator and \( \eta \) is its first PF descendent. The closed algebra generated by the currents \( T(z) \) and \( G(z) \) is the level \( K \) FSCA (Since \( \epsilon \) does not exist for \( K = 1 \), \( G(z) \) is absent in this case and we recover the Virasoro algebra). The central charge of this algebra is \( c_0 \). If we turn on the background charge of the \( \varphi \) boson, the algebra, generated by the appropriately modified \( T(z) \) and \( G(z) \) currents, remains closed. We note that there exist other FSCAs with additional fractional supercurrents [6].

The following OPEs define the level \( K \) FSCA with arbitrary central charge \( c \):
\[
T(z)T(w) = \frac{(c/2)}{(z-w)^4} + \frac{2T(w)}{(z-w)^2} + \frac{\partial T(w)}{(z-w)} + ... ,
\]
\[
T(z)G(w) = \frac{\Delta}{(z-w)^2} G(w) + \frac{1}{(z-w)} \partial G(w) + ... ,
\]
\[
G(z)G(w) = (z-w)^{-2\Delta} \left\{ \frac{c}{\Delta} + 2(z-w)^2 T(w) + ... \right\} + \lambda(c)(z-w)^{-\Delta} \left\{ G(w) + \frac{1}{2}(z-w) \partial G(w) + ... \right\}.
\]

The associativity condition for this algebra fixes the structure constant \( \lambda^2(c) \):
\[
\lambda^2(c) = \frac{2K^2 c_{111}^2}{3(K + 4)^2} (c_1 - c) .
\]

Here \( c_{111} \) is the \( SU(2)_K \) structure constant for the OPE of two chiral spin-1 primary fields to give a chiral spin-1 field and [6]
\[
c_1 \equiv 24/K + c_0 .
\]

As mentioned earlier, for \( K = 1 \) the \( G \) current is absent. For \( K = 2, c_{111} = 0, \lambda(c) = 0 \) and \( G \) is the usual supercurrent. It is a hermitian current of conformal dimension 3/2. For \( K \geq 3, c_{111} \) is non-vanishing and
\[ G^\dagger = \alpha G \ , \quad (2.14) \]

where

\[
\alpha = \begin{cases} 
+1 & \text{if } c \leq c_1 , \\
-1 & \text{otherwise .} 
\end{cases} \quad (2.15)
\]

For definiteness we will take the branch of Eqn.\((2.12)\) that is positive for \(c < c_1\).

The vertex operators \(\chi^j_m(0)\), labeled by the \(SU(2)_K\) quantum numbers \(j\) and \(m\), create states in modules of the level \(K\) FSCA: \(|\chi^j_m\rangle \equiv \chi^j_m(0)|0\rangle\). The fields primary with respect to \(G(z)\) and \(T(z)\) have the form \(\chi^j_m(z)\) (\(\chi^j_{-j}(z) = (\chi^j_m)^\dagger\) are also primary). \(G(z)\) has \(j = 1\) and \(m = 0\) and in terms of modings acts on \(|\chi^j_m\rangle\) in the same manner as \(\epsilon\) acts on the PF fields \(\phi_m^j\):

\[
G(z)\chi^j_m(0) = \sum_n \left[ z^{n+2j/(K+2)} G_{n-2(j+1)/(K+2)} + z^{n-2/(K+2)} G_{n} + z^{n-2(j+1)/(K+2)} G_{n+2j/(K+2)-1} \right] \chi^j_m(0) , \quad (2.16)
\]

where the resulting states have spins \(j + 1, j\) and \(j - 1\), respectively (The terms not allowed by the fusion rules \((2.14)\) are absent). The positive modes of the \(T\) and \(G\) currents annihilate the highest weight states created by the primary fields:

\[
L_0|h\rangle = h|h\rangle , \quad L_n|h\rangle = G_r|h\rangle = 0 , \quad n, r > 0 . \quad (2.17)
\]

Here \(r\) is in general a rational number.

The module \([\chi^j_j]\) built from the highest weight state \(|\chi^j_j\rangle\) is spanned by the states

\[
L_{-n_1}^{\lambda_1} \ldots L_{-n_p}^{\lambda_p} G_{-r_1}^{\rho_1} \ldots G_{-r_q}^{\rho_q} |\chi^j_j\rangle \ , \quad (2.18)
\]

where the ordering is the same as for the states

\[
a_{-n_1}^{\lambda_1} \ldots a_{-n_p}^{\lambda_p} \epsilon_{-r_1}^{\rho_1} \ldots \epsilon_{-r_q}^{\rho_q} \phi^j_j(0) |0\rangle , \quad n_i > 0 , \quad r_k \geq 0 . \quad (2.19)
\]

Here \(a_{-n_i}\) represent creation modes of the \(\varphi\) boson, while \(\epsilon_{-r_k}\) stand for modes of the \(Z_K\) PF energy operator. Note that \(\lambda_i\) can be any non-negative integers, while \(\rho_k\) can take only two values: zero and unity.

The naive counting of states \(|\chi^j_m\rangle\) of the form \((2.18)\) is identical to the counting of states \((2.19)\) with the same quantum numbers \(j\) and \(m\) and is given by the level \(K\) string functions \(C^{2j}_{2m}\) \((2.19)\)

\[
C^{2j}_{2m} = \frac{q^{d_m}}{\eta^3} \sum_{r,s=0}^{\infty} (-1)^{r+s} q^{s(s+1)/2+r(r+1)/2+s(r+1)} \\
\times \left( q^{s(j-m)+r(j+m)} - q^{K+1-2j+m[K+1-j+m]+r[K+1-j-m]} \right) . \quad (2.20)
\]

Here \(\eta\) is the Dedekind \(\eta\)-function

\[
\eta(q) = q^{1/24} \prod_{n=1}^{\infty} (1 - q^n) , \quad (2.21)
\]
and
\[ d^j_m = \Delta^j_m + \frac{1}{4(K + 2)}. \] (2.22)

For later convenience we define the level of a state \((2.18)\) as
\[ N = N_0 + \Delta^j_j, \] (2.23)
where \(\Delta^j_j\) is given by \((2.4)\) and
\[ N_0 = \sum_{i=1}^p \lambda_i n_i + \sum_{k=1}^q \rho_k r_k \] (2.24)
is (generically a rational number) the level of the state above the highest weight state \(|\chi^j_j\rangle\).
The number of its descendents \(i.e.,\) the number of states in the module \([\chi^j_j]\) at level \(N\) is given by \(P_j(N)\), the latter being a coefficient in the power expansion
\[ C_{2m}(q) \equiv \sum_j C_{2m}^j(q) = q^{-\frac{2m}{24}} \sum P_m(N)q^N. \] (2.25)
The sum over \(j\) includes its values allowed by the fusion rules \((2.6)\) up to identifications \((2.5)\), \(i.e.,\) each independent string function appears once and only once.

We note that the \(G\) current modings when acting on the highest weight states \(|\chi^j_j\rangle\) and \(|\chi^{K/2-j}_j\rangle\) \((0 < j < K/4)\) are the same. Hence the modules \([\chi^j_j]\) and \([\chi^{K/2-j}_j]\) are structurally identical.

We end this subsection with the following remark. An important feature of level \(K \geq 3\) FSCAs is the appearance of cuts in the \(GG\) OPE \((2.11)\). Since there are two different cuts on the right hand side, upon continuation of a correlation function involving \(G(z)G(w)\) OPE along a contour interchanging \(z\) and \(w\) in the complex plane, it is inconsistent for the correlator to pick up a simple phase; that is, the current \(G\) is non-abelianly braided. There is only one FSCA, namely, the spin-4/3 \((K = 4)\) parafermion current algebra, where it is possible to split the fractional supercurrent so that the split currents are abelianly braided. This makes for important technical simplifications. For example, it enables one to write down the generalized commutation relations between the modes of the split currents and to explicitly calculate the Kac and new determinants. For other FSCAs the generalized commutation relations have not yet been written down because of technical difficulties arising from their non-abelian braiding properties. Hence, the spin-4/3 parafermion current algebra provides an important check on the Kac and new determinant formulae presented in the next subsection, and we consider it in detail in subsection C.

B. Kac and New Determinant Formulae

As mentioned earlier, so far the Kac determinant formulae have been known only for the Virasoro and superconformal algebras. These Kac determinants are polynomials in the highest weight \(h\). One expects the determinants of inner products of states in modules of FCSAs also be polynomials in \(h\). This is the case for some modules, but not for all.
There are three different types of modules in the Fock space of the level $K$ FSCA.

(i) In the module $|\chi_0^0\rangle$, built from the highest weight state $|\chi_0^0\rangle$, the counting of states is given by $C_0$ defined in (2.25). The fractional supercurrent zero mode $G_0$ does not act on $|\chi_0^0\rangle$ and the determinant of inner products of states at any level is a polynomial in $h$.

(ii) For $0 < j \neq K/4$ the highest weight state $|\chi_j^j\rangle$ is an eigenstate of $G_0$. The counting of states in the module $|\chi_j^j\rangle$ given by $C_{2j}$. The determinant of inner products of states at any level is a polynomial in the $G_0$ eigenvalue, but is not a polynomial in $h$. However, as mentioned earlier, the modules $|\chi_j^j\rangle$ and $|\chi_j^{K/2-j}\rangle$ ($0 < j < K/4$) are structurally identical.

This allows us to construct the highest weight state $|h\rangle$ (as a linear combination of $|\chi_j^j\rangle$ and $|\chi_j^{K/2-j}\rangle$) that is no longer an eigenstate of $G_0$. The determinant of inner products of states in the module $|\chi_j^j \oplus \chi_j^{K/2-j}\rangle$ built from $|h\rangle$ is a polynomial in $h$. The number of states is given by $C_{2j} + C_{K-2j}(\equiv 2C_{2j})$.

(iii) If $K$ is even, there is the third type of modules, namely, the $|\chi_{K/2}^{K/2}\rangle$ module. $G_0$ is not necessarily diagonal with respect to the highest weight state $|\chi_{K/4}^{K/4}\rangle$. However, the determinant of inner products of states is always a polynomial in $h$. The counting of states is given by $C_{K/2}$.

Now we define $\det(M^{(j)}(N))$ as the determinant of inner products of states (2.18) at level $N$. It is a generalization of the Kac determinants for the Virasoro and superconformal algebras. For the modules $|\chi_0^0\rangle$, $|\chi_{K/4}^{K/4}\rangle$ and $|\chi_j^j \oplus \chi_j^{K/2-j}\rangle$ ($0 < j < K/4$) the determinant is a polynomial in $h$, and we will refer to it as the Kac determinant. For the module $|\chi_j^j\rangle$ ($0 < j \neq K/4$) the determinant is a polynomial in the $G_0$ eigenvalue, but is not a polynomial in $h$. We will refer to these determinants as the new determinants. The Kac determinant for the module $|\chi_j^j \oplus \chi_j^{K/2-j}\rangle$ factorizes into a product of the new determinants for the $|\chi_j^j\rangle$ and $|\chi_j^{K/2-j}\rangle$ modules.

First we present (up to positive normalization constants independent of the highest weight $h$ and the central charge $c$) the Kac and new determinant formulae for the modules described above. Then we give explicit formulae for the zeros of the Kac and new determinants. Finally, we comment on some features of the determinants for FSCAs.

• (i) The module $|\chi_0^0\rangle$.

The Kac determinant reads

$$\det(M^{(0)}(N_0)) = \alpha^{Q(0)} \prod_{r,s} [\alpha(h - h_{r,s})]^{P_k(N_0 - rs/K)},$$

where $r, s \in \mathbb{N}$, $s - r = 0 \mod K$ and $rs/K \leq N_0$; $2\ell = (s + r) \mod K$ and

$$Q(0) = \sum_n P_0(N_0 - n^2),$$

where $n \in \mathbb{N}$, $n^2 \leq N_0$.

• (ii) The module $|\chi_{K/4}^{K/4}\rangle$, $K$ is even.

(a) $G_0$ is diagonal with respect to the highest weight state $|h; g^\pm\rangle$:

$$G_0|h; g^\pm\rangle = g^\pm|h; g^\pm\rangle.$$

With the appropriate normalization of $G_0$, the eigenvalues $g^\pm$ are given by
\[ g^\pm = \pm \sqrt{h - c/24} . \]  

The Kac determinants for the modules built from \( |h; g^\pm\rangle \) read

\[
\text{det}(\mathcal{M}_{(\pm)}^{(K/4)}(N)) = \alpha^{Q(K/4)} \prod_{r,s}(h - h_{r,s})^{P_r(N-rs/K)} ,
\]

where \( N \equiv N_0 + K/8(K+2) \), \( s - r = K/2 \) mod \( K \) and \( rs/K \leq N \); \( 2\ell = (s + r) \) mod \( K \) and

\[
Q(K/4) = P_{K/4}(N)/2 + \sum_n P_{K/4}(N - n^2) .
\]

Here \( n^2 \leq N \).

(b) \( G_0 \) is not diagonal with respect to the highest weight state \( |h\rangle \).

The Kac determinant reads

\[
\text{det}(\mathcal{M}^{(K/4)}(N)) = \alpha^{Q(K/4)}(h - c/24)^{P_{K/4}(N)/2} \prod_{r,s}(h - h_{r,s})^{P_r(N-rs/K)} ,
\]

where \( r, s, \ell \) and \( N \) are the same as in Eqn.(2.30).

(iii) The module \([\chi_j^2]\), \( 0 < j \neq K/4 \).

(a) The highest weight state \( |\chi_j^2\rangle \equiv |h; g^j\rangle \) is an eigenstate of \( G_0 \):

\[
G_0 |h; g^j\rangle = g^j |h; g^j\rangle .
\]

With the appropriate normalization of \( G_0 \), the eigenvalue \( g^j \) is given by

\[
g^j = a^j \sqrt{c_1 - c + \text{sgn}(j - K/4) \sqrt{h - h_0^j}} ,
\]

where

\[
a^j = \frac{|K - 4j|}{\sqrt{24(K + 4)}} ,
\]

and we define for arbitrary \( j \)

\[
h_0^j \equiv \frac{c - c_0}{24} + \Delta_j^j = \frac{c}{24} - \frac{(K - 4j)^2}{8K(K + 2)} .
\]

The new determinant reads

\[
\text{det}(\mathcal{M}^{(j)}(N)) = \alpha^{Q(j)} \prod_{r,s}[(\alpha(g^j - g_{r,s})(g^j - g_{s,r})]^{P_r(N-rs/K)} ,
\]

where \( N = N_0 + \Delta_j^j \), \( s - r = 2j \) mod \( K \) and \( rs/K \leq N \); \( 2\ell = (s + r) \) mod \( K \) and

\[
Q(j) = \sum_n P_j(N - n^2) .
\]

(b) The module \([\chi_j^2 \oplus \chi_{K/2-j}^2]\), \( 0 < j < K/4 \).
The highest weight state \(|h\rangle\) is a linear combination of \(|\chi^j_j\rangle\) and \(|\chi_{K/2-j}^{K/2-j}\rangle\) such that \(G_0\) is not diagonal with respect to \(|h\rangle\).

The Kac determinant reads

\[
\det(\mathcal{M}^{(j)}_0(N)) = [\alpha(h - h_0^j)]^{P_j(N)} \prod_{r,s} [(h - h_{r,s})(h - h_{s,r})]^{P_{r(N-rs/K)}} ,
\]

where \(r, s, \ell\) are the same as in (2.37).

The Kac determinant (2.39) factorizes as follows:

\[
\det(\mathcal{M}^{(j)}_0(N)) = [\alpha(h - h_0^j)]^{P_j(N)} \det(\mathcal{M}^{(j)}(N)) \det(\mathcal{M}^{(K/2-j)}(N)) .
\]

The zeros of the Kac determinants (2.26), (2.30), (2.32) and (2.39) read

\[
h_{r,s} = \frac{c - c_0}{24} + \Delta^j_0 + \frac{1}{96} \left( (r + s) \sqrt{c_0 - c} + (r - s) \sqrt{c_1 - c} \right)^2 = h_0^j + b_{r,s}^2 ,
\]

where

\[
b_{r,s} = \frac{1}{\sqrt{96}} \left( (r + s) \sqrt{c_0 - c} + (r - s) \sqrt{c_1 - c} \right) .
\]

The zeros of the new determinants (2.37) are given by

\[
g_{r,s} = a^j \sqrt{c_1 - c} + \text{sgn}(j - K/4) b_{r,s} .
\]

Here \(2j = (s - r) \mod K\). The relation between the zeros of the Kac and corresponding new determinants reads

\[
(g^j - g_{r,s})(g^{K/2-j} - g_{r,s}) = -(h - h_{r,s}) .
\]

Their orders \(P_m(N)\) are given by

\[
C_{2m}(q) \equiv \sum_j C_{2m}^j(q) = q^{-\frac{K}{8}} \sum_{m} P_m(N) q^N .
\]

The sum over \(j\) includes its values allowed by the fusion rules (2.6) up to identifications (2.5), i.e., each independent string function appears once and only once.

We conclude this subsection with the following remarks:

(a) We define the Kac and new determinants for a linearly independent set of states where the total number of the \(T\) and \(G\) modes is minimal. If this number is not minimal, the determinant will be different by an overall factor (that can be either positive or negative) independent of \(h\), but in general dependent on \(c\).

(b) For the Virasoro algebra the \(G\) current is absent, while for the superconformal algebra it is always hermitian. Therefore, for the \(K = 1, 2\) cases \(\alpha \equiv 1\) in all of the above formulae. For level \(K \geq 3\) FSCAs, \(\alpha\) flips sign as the central charge exceeds \(c_1\) defined in (2.13). Hence in the Kac and new determinants there appears the factor of \(\alpha\) to the power (defined mod 2) determined by the total number of the \(G\) current modes in a linearly independent set of states with minimal counting.

(c) For level \(K \geq 3\) FSCAs, the orders of some zeros of the Kac and new determinants for a given module come from the string functions that give the counting of states in a different module. This is never the case for the Virasoro and superconformal algebras.

The examples in the next subsection illuminate these issues.
C. Examples

Now we consider some concrete cases in order to clarify our conventions and notations. We give the Kac and new determinant formulae up to normalization constants independent of the highest weight $h$ and the central charge $c$.

- (i) The Virasoro algebra ($K = 1$).

We have only one module $[\chi_0^0]$ and the counting of states is given by the string function

$$ C_0 \equiv C_0^0 = 1/\eta(q) = q^{-1/24} \prod_{n=1}^{\infty} (1 - q^n)^{-1} . $$ (2.46)

Eqn.(2.26) (with $\alpha \equiv 1$) reduces to the well-known Kac determinant formula for the Virasoro algebra.

- (ii) The superconformal algebra ($K = 2$).

(a) The Neveu-Schwarz module, or the $[\chi_0^0]$ module. According to Eqn.(2.25) the counting of states is given by

$$ C_0 \equiv C_0^0 + C_0^2 = q^{-1/16} \prod_{n=1}^{\infty} \frac{1 + q^{n-1/2}}{1 - q^n} , $$ (2.47)

Eqn.(2.26) (again with $\alpha \equiv 1$) now reduces to the Kac determinant formula for the Neveu-Schwarz module.

(b) The Ramond module, or the $[\chi_1^{1/2}]$ module. This is the simplest $[\chi_{K/4}^{K/4}]$ module. The counting is given by the string function

$$ C_1 \equiv C_1^1 = q^{-1/16} \left(q^{1/16} \prod_{n=1}^{\infty} \frac{1 + q^n}{1 - q^n}\right) , $$ (2.48)

The level of a state in this module according to Eqn.(2.23) is defined as $N = N_0 + 1/16$, where $N_0$ is the level of the state above the highest weight state. Eqn.(2.32) is then the Kac determinant formula for the Ramond module.

To illuminate the Kac determinant formula (2.30) for the $[\chi_{K/4}^{K/4}]$ module, we consider the Ramond module built from an eigenstate $|h; g^\pm\rangle$ of the supercurrent zero mode $G_0$. Due to the zero mode relation

$$ G_0^2 = L_0 - c/24 , $$ (2.49)

we have

$$ g^\pm = \pm \sqrt{h - c/24} . $$ (2.50)

A linearly independent set of states at level $17/16 = 1 + 1/16$ reads

$$ |\phi_1^\pm\rangle = L_{-1}|h; g^\pm\rangle , \quad |\phi_2^\pm\rangle = G_{-1}|h; g^\pm\rangle . $$ (2.51)

The Kac determinants of these states can be computed using the superconformal mode algebra:
\[
\det(M(\pm)) = 4(h - h_{1,2})(h - h_{2,1}),
\] (2.52)

where \(h_{1,2}\) and \(h_{2,1}\) are given by Eqn. (2.41) at \(K = 2\) and \(j = 1/2\). Note that at this level the Kac determinant is a polynomial in \(h\). This holds at all levels for all the \([\chi_{K/4}]\) modules.

- (iii) The spin-4/3 parafermion current algebra \((K = 4)\).

The chiral fractional supercurrent \(G(z)\), whose conformal dimension is \(4/3\), can be split into two pieces:

\[
G(z) = \begin{cases} 
G^+(z) + G^-(z) & \text{if } c \leq 8, \\
(i(G^-(z) - G^+(z)) & \text{otherwise}.
\end{cases}
\] (2.53)

The currents \(G^\pm(z)\) are abelianly braided (i.e., parafermionic), and with the \(T(z)\) current form the closed spin-4/3 parafermion current algebra of Zamolodchikov and Fateev [8].

\[
G^\pm(z)G^\pm(w) = \frac{\lambda^\pm}{(z-w)^{4/3}} \left\{ G^\mp(w) + \frac{1}{2} (z-w) \partial G^\mp(w) + \ldots \right\},
\]

\[
G^+(z)G^-(w) = \frac{1}{(z-w)^{8/3}} \left\{ \frac{3c}{8} + (z-w)^2 T(w) + \ldots \right\},
\] (2.54)

\[
T(z)G^\pm(w) = \frac{(4/3) G^\pm(w)}{(z-w)^2} + \partial G^\pm(w) + \ldots,
\]

\[
T(z)T(w) = \frac{(c/2)}{(z-w)^4} + \frac{2T(w)}{(z-w)^2} + \partial T(w) + \ldots,
\]

where

\[
\lambda^\pm = \alpha \lambda^- = \lambda = \sqrt{\frac{8-c}{6}}, \quad \alpha = \begin{cases} +1 & \text{if } c \leq 8, \\
-1 & \text{otherwise}.
\end{cases}
\] (2.55)

The \(G^\pm\) currents have the following hermiticity properties:

\[
(G^\pm)^\dagger = \alpha G^\mp.
\] (2.56)

The algebra (2.54) obeys a \(\mathbb{Z}_3\) symmetry. The \(G^\pm(z)\) and \(T(z)\) currents have \(\mathbb{Z}_3\) charges \(q = \pm 1\) and zero, respectively. The \(\mathbb{Z}_3\) charge \(q\) is defined mod 3.

The modes of the currents

\[
T(z) = \sum_n z^{-n-2} L_n, \quad G^\pm(z) = \sum_r z^{-r-4/3} G^\pm_r
\] (2.57)

satisfy the commutation relations

\[
[L_m, L_n] = (m-n)L_{m+n} + \frac{c}{12} (m^3-m) \delta_{m+n},
\]

\[
[L_m, G^\pm_r] = \left( \frac{m}{3} - r \right) G^\pm_{m+r}.
\] (2.58)

The modings of the currents \(G^\pm(z)\) and the generalized commutation relations (GCRs) for them depend on a representation of the algebra. First we consider the representations
that correspond to the integer spin \( j \) modules. The Fock space falls into sectors \( \mathcal{H}_q \) labeled by their \( \mathbb{Z}_3 \) charge. The currents \( G^\pm \) act on the Fock space sectors according to the rules

\[
G^\pm : \mathcal{H}_q \to \mathcal{H}_{q \pm 1} .
\]

With these actions, the mode expansions of the \( G^\pm \) currents are defined as

\[
G^\pm(z)\chi_q(0) = \sum_n z^{n+q/3} G^\pm_{-n-(1+q)/3} \chi_q(0) ,
\]

where \( \chi_q \) is an arbitrary state in \( \mathcal{H}_q \). These mode expansions can be inverted to give

\[
G^\pm_{-n-(1+q)/3} \chi_q(0) = \oint_\gamma \frac{dz}{2\pi i} z^{-q/3} G^\pm(z) \chi_q(0) .
\]

Here, \( \gamma \) is a contour encircling the origin once, where \( \chi_q(0) \) is inserted.

The \( G_r \) modes satisfy the following GCRs

\[
\sum_{\ell=0}^{\infty} c^{(1/3)}_\ell \left[ G_{\ell+q/3}^\pm G_{\ell+q/3+1}^\mp + G_{\ell+q/3+1}^\pm G_{\ell+q/3}^\mp \right] = \lambda^\pm G_{\ell+q/3+n+m}^\mp ,
\]

\[
\sum_{\ell=0}^{\infty} c^{(-1/3)}_\ell \left[ G_{\ell+q/3}^\pm G_{\ell+q/3-1}^\mp + G_{\ell+q/3-1}^\pm G_{\ell+q/3}^\mp \right] =
\]

\[
L_{n+m} + \frac{3c}{16} \left( n + 1 + \frac{q}{3} \right) \left( n + \frac{q}{3} \right) \delta_{n+m} ,
\]

when acting on a state with the \( \mathbb{Z}_3 \) charge \( q \). Here \( c^{(a)}_\ell \) are the binomial coefficients

\[
(1-x)^a = \sum_{\ell=0}^{\infty} c^{(a)}_\ell x^\ell .
\]

According to the hermiticity assignments (2.56) and the mode expansions (2.60), the hermiticity properties of the \( G_r^\pm \) modes are \( (G_r^+)^\dagger = \alpha G_{-r}^\mp \).

The highest weight states \( |h; q\rangle \) with the \( \mathbb{Z}_3 \) charge \( q \) satisfy the following conditions

\[
L_0 |h; q\rangle = h |h; q\rangle , \quad L_n |h; q\rangle = 0 , \quad G_{n-\frac{1+q}{3}}^\pm |h; q\rangle = 0 , \quad n > 0 .
\]

Next we consider the representations of the algebra corresponding to the half odd integer spin \( j \) modules. The fields \( z^{1/3} G^\pm(z) \chi_R(0) \) (where \( \chi_R \) is an arbitrary state in the Fock space) are double-valued analytic functions, \( i.e., \) the fields \( \chi_R(z) \) create square root cuts in the complex plane. For the purpose of calculating the Kac and new determinants it suffices to consider only the states \( \chi_R(0) \) that are \( \mathbb{Z}_3 \) singlets. The modings of the \( G^\pm \) currents then are given by the following OPEs

\[
G^\pm(z) \chi_R(0) = 2^{-\frac{2}{3}} \sum_{n=-\infty}^{\infty} z^{n/2-4/3} G_{-n/2}^\pm \chi_R(0) , \quad n \in \mathbb{Z} ,
\]

where the overall normalization is chosen for later convenience. These mode expansions can be inverted to give
\[ G^{\pm}_{n/2} \chi_R(0) = 2^{-1/4} \frac{1}{2 \pi i} \oint_{\gamma} z^{1/3+n/2} G^{\pm}(z) \chi_R(0). \quad (2.67) \]

Here, \( \gamma \) is a contour encircling the origin twice, where \( \chi_R(0) \) is inserted.

The modes \( G^{\pm}_{n/2} \) are related to each other via \( G^+_{n/2} = (-1)^n \alpha G^-_{n/2} \) due to the single-bypass relations discussed in Ref [8]. Defining \( \tilde{G}_{n/2} \equiv G^-_{n/2} \), we find that the modes \( \tilde{G}_{n/2} \) satisfy the following GCR

\[ \sum_{\ell=0}^{\infty} D^{(\ell)}_{(\alpha, \beta)} \left[ \tilde{G}_{m-\ell} \tilde{G}_{m+\ell} + \tilde{G}_{m+\ell} \tilde{G}_{m-\ell} \right] = \frac{\lambda}{2} (-1)^{n+m} \tilde{G}_{n+m} + \alpha \left[ (-1)^n + (-1)^m \right] \left\{ L_{n+m} + \frac{3c}{8} \left( \frac{n^2}{8} - \frac{5}{48} \right) \delta_{n+m} \right\}, \quad (2.68) \]

where \( D^{(\ell)}_{(\alpha, \beta)} \) are the binomial coefficients

\[ (1-x)^\alpha (1+x)^\beta = \sum_{\ell=0}^{\infty} D^{(\ell)}_{(\alpha, \beta)} x^\ell. \quad (2.69) \]

According to the hermiticity assignments \( (2.56) \) and the mode expansions \( (2.66) \), the hermiticity properties of the \( \tilde{G}_{n/2} \) modes satisfy the following conditions

\[ L_0 |h\rangle_R = h |h\rangle_R, \quad L_n |h\rangle_R = 0, \quad \tilde{G}_{n/2} |h\rangle_R = 0, \quad n > 0. \quad (2.70) \]

Now we turn to the description of the spin-4/3 parafermion current algebra modules where the counting of states is given by the \( K = 4 \) string functions [16,17]

\[ C_0^0 + C_0^4 = \frac{1}{\eta^2(q)} \left( \sum_{n=-\infty}^{\infty} q^{3n^2} \right), \]
\[ C_2^2 = \frac{q^{1/12}}{\eta^2(q)} \left( \sum_{n=-\infty}^{\infty} q^{3n^2+n} \right), \]
\[ C_0^4 = \frac{q^{1/3}}{\eta^2(q)} \left( \sum_{n=-\infty}^{\infty} q^{3n^2+2n} \right), \]
\[ C_2^4 = C_2^0 = \frac{q^{3/4}}{\eta^2(q)} \left( \sum_{n=0}^{\infty} q^{3n^2+3n} \right), \]
\[ C_1^1 + C_1^3 = C_3^3 + C_3^1 = 1/\eta(q^{1/2}). \quad (2.71) \]

According to Eqn.\((2.25)\) we define

\[ C_0 \equiv C_0^0 + C_0^2 + C_0^4 = q^{-1/12} \sum P_0(N) q^N = q^{-1/12} (1 + q^{1/3} + \ldots), \]
\[ C_2 \equiv C_0^2 + C_2^2 + C_2^4 = q^{-1/12} \sum P_1(N) q^N = q^{-1/12} (q^{1/12} + 2q^{3/4} + \ldots), \]
\[ C_1 \equiv C_1^1 + C_1^3 = q^{-1/12} \sum P_{1/2}(N) q^N = q^{-1/12} (q^{1/16} + q^{9/16} + \ldots), \]
\[ C_3 \equiv C_3^1 + C_3^3 = q^{-1/12} \sum P_{3/2}(N) q^N = C_1. \quad (2.72) \]
We also define the following combination of the string functions:

$$C_0 + C_2 = q^{-1/12} \sum P(N)q^N.$$

(2.73)

(a) The S-module, or the $[\chi_0^0]$ module. This module is built from the highest weight state $|h; 0\rangle$. We define the level of a state as $N = N_0$, where $N_0$ is the level of the state above $|h; 0\rangle$. Eqn. (2.60) gives $N \in \mathbb{Z}$ or $N \in \mathbb{Z} + 1/3$. The $S$-module falls into three submodules, $S^{(0)}$ and $S^{(\pm)}$, labeled according to the $\mathbb{Z}_3$ charge of the states. The $S^{(0)}$-submodule contains the states (with quantum numbers $j = 0, 2$ and $m = 0$) at integer levels, while the $S^{(\pm)}$-submodules consist of the states ($j = 1$ and $m = 0$) at levels $N \in \mathbb{Z} + 1/3$. The number of states at level $N$ in each submodule is given by $P_0(N)$ defined in (2.72).

For the $S$-module we define the Kac determinant at integer levels as that of the $S^{(0)}$-submodule, and at levels $N \in \mathbb{Z} + 1/3$ as that of the $S^{(\pm)}$-submodule. The Kac determinants for the $S^{(+)}$- and $S^{(-)}$-submodules are identical. We choose the normalization $\langle h; 0 | h; 0 \rangle = 1$.

The Kac determinant formula for the $S$-module is given by

$$\det(M^S_N) = \alpha \frac{Q(S)}{\prod_{r,s} (h - h^{S}_{r,s})^{P(N - rs/4)}}.$$

(2.74)

where $r, s \in \mathbb{N}$, $s - r = 0 \mod 4$, $rs/4 \leq N$, and

$$Q(S) = \sum_n P_0(N - n^2/4).$$

(2.75)

Here $n \in \mathbb{N}$ and $n^2/4 \leq N$. The zeros $h^{S}_{r,s}$ of the Kac determinant (2.74) are given by Eqn. (2.41)

$$h^{S}_{r,s} = \frac{c - 2}{24} + \frac{1}{96} \left((r + s)\sqrt{2 - c} + (r - s)\sqrt{8 - c}\right)^2.$$

(2.76)

Note that the order of the zeros $P(N - rs/4)$ comes from the sum $C_0 + C_2$ defined in Eqn. (2.73), not just from $C_0$.

We present a few low lying levels explicitly to clarify the notation. There is $P_0(1/3) = 1$ state in the $S^{(\pm)}$-submodule at level 1/3:

$$|\phi\rangle = G^{-}_{-1/3}|h; 0\rangle.$$

(2.77)

All the other states at this level are linearly dependent on (2.77). For instance, using the GCR (2.62) we have $G_0^+ G^{+}_{-1/3}|h; 0\rangle = (\lambda^+/2) G^{-}_{-1/3}|h; 0\rangle$, and we choose the state (2.77), since it has the minimal number of the $G$ creation operators. The Kac determinant can be calculated using (2.58) and (2.63):

$$\det(M^S_{1/3}) = \alpha (h - h_{1,1}).$$

(2.78)

At level 1 there are $P_0(1) = 2$ states in the $S^{(0)}$-submodule

$$|\phi_1\rangle = L_{-1}|h; 0\rangle, \quad |\phi_2\rangle = G^{-}_{-1} G^{+}_{-1/3}|h; 0\rangle.$$

(2.79)

The Kac determinant can be obtained using the commutation relations (2.58) and GCRs (2.62) and (2.63):
\[ \det(\mathcal{M}_1^S) = \frac{4}{3}(h - h_{1,1})^2(h - h_{2,2}) . \] (2.80)

At the next level \( N = 4/3 \) there are \( P_0(4/3) = 3 \) states
\[ |\phi_1\rangle = L_{-1}G_{-\frac{1}{3}}^-|h; 0\rangle , \quad |\phi_2\rangle = G_{-1}^+G_{-\frac{1}{3}}^+|h; 0\rangle , \quad |\phi_3\rangle = G_{-\frac{1}{3}}^-|h; 0\rangle , \] (2.81)
and the Kac determinant is given by
\[ \det(\mathcal{M}_{4/3}^S) = \frac{4}{3}(h - h_{1,1})^2(h - h_{2,2})(h - h_{1,5})(h - h_{5,1}) . \] (2.82)

(b) The \( D \)-module, or the \([\chi_1^1]\) module. This module (of the \([\chi_{K/4}^K]\) type) is built from the highest weight state \(|h; -1\rangle\). We define the level of a state as \( N = N_0 + 1/12 \), where \( N_0 \)

is the level of the state above \(|h; -1\rangle\). Eqn. (2.60) gives \( N \in \mathbb{Z} + 1/12 \) or \( N \in \mathbb{Z} + 3/4 \). The \( D \)-module falls into three submodules, \( D^{(0)} \) and \( D^{(\pm)} \). The \( D^{(\pm)} \)-submodules consist of the states \((j = 1 \text{ and } m = 1)\) at levels \( N \in \mathbb{Z} + 1/12 \), while the \( D^{(0)} \)-submodule contains the states \((j = 0, 2 \text{ and } m = 1)\) at levels \( N \in \mathbb{Z} + 3/4 \). The number of states at level \( N \) in each submodule is given by \( P_1(N) \) defined in (2.72).

For the \( D \)-module we define the Kac determinant at levels \( N \in \mathbb{Z} + 1/12 \) as that of the \( D^{(-)} \)-submodule, and at levels \( N \in \mathbb{Z} + 3/4 \) as that of the \( D^{(0)} \)-submodule. The Kac determinants for the \( D^{(+)} \)- and \( D^{(-)} \)-submodules are identical for \( N > 1/12 \). We choose the standard normalization \( \langle h; -1|h; -1\rangle = 1, G_0^\dagger|h; -1\rangle = \sqrt{h - c/24}|h; +1\rangle \).

The Kac determinant formula for the \( D \)-module is given by
\[ \det(\mathcal{M}_N^D) = \alpha^{Q(D)}(h - c/24)^{P_1(N)/2}\prod_{r,s}(h - h_{r,s}^D)^{P(N - rs/4)} . \] (2.83)

where \( s - r = 2 \text{ mod } 4, rs/4 \leq N > 1/12 \) and
\[ Q(D) = \frac{P_1(N)/2 + \sum_n P_1(N - n^2)}{2} . \] (2.84)

Here \( n^2 \leq N \). The zeros \( h_{r,s}^D \) are given by
\[ h_{r,s}^D = \frac{c}{24} + \frac{1}{96}\left((r + s)\sqrt{2 - c} + (r - s)\sqrt{8 - c}\right)^2 . \] (2.85)

For example, at level \( 3/4 = 2/3 + 1/12 \) there are \( P_1(3/4) = 2 \) states in the \( D^{(0)} \)-submodule:
\[ |\phi_1\rangle = G_{-\frac{1}{4}}^+|h; -1\rangle , \quad |\phi_2\rangle = G_{-\frac{1}{4}}^-G_0^\dagger|h; -1\rangle , \] (2.86)
and
\[ \det(\mathcal{M}_{3/4}^D) = \alpha(h - c/24)(h - h_{1,3})(h - h_{3,1}) . \] (2.87)

At level \( 13/12 = 1 + 1/12 \) there are \( P_1(13/12) = 2 \) states in the \( D^{(-)} \)-submodule
\[ |\phi_1\rangle = L_{-1}|h; -1\rangle , \quad |\phi_2\rangle = G_{-\frac{1}{1}}^+G_0^\dagger|h; -1\rangle . \] (2.88)
and the Kac determinant reads

\[
\det(M_{13/12}^D) = \frac{4}{3} (h - c/24)(h - h_{1,3})(h - h_{3,1}).
\] (2.89)

We also give the first non-trivial level in the \(D\)-module built from the highest weight states that are the eigenstates of the fractional supercurrent zero mode \(G_0\):

\[
G_0|h; g^\pm\rangle = g^\pm|h; g^\pm\rangle.
\] (2.90)

With our normalization of \(G_0\) in Eqn.(2.29),

\[
G_0 = \begin{cases} 
G_0^+ + G_0^- & \text{if } c \leq 8 \\
i(G_0^- - G_0^+) & \text{otherwise}.
\end{cases}
\] (2.91)

The eigenstates of \(G_0\) read

\[
|h; g^\pm\rangle = \begin{cases} 
\frac{1}{\sqrt{2}} (|h; -1\rangle \pm |h; 1\rangle) & \text{if } c \leq 8 \\
\frac{-1}{\sqrt{2}} (|h; -1\rangle \pm i|h; 1\rangle) & \text{otherwise}.
\end{cases}
\] (2.92)

They correspond to the eigenvalues \(g^\pm = \pm \sqrt{h - c/24}\). At level 3/4 we have the following states

\[
|\phi_1^\pm\rangle = G_{-3/4}^+ |h; g^\pm\rangle, \quad |\phi_2^\pm\rangle = G_{-3/4}^- |h; g^\pm\rangle.
\] (2.93)

The corresponding Kac determinants read

\[
\det(M_{3/4}^{(\pm)}) = \alpha(h - h_{1,3})(h - h_{3,1}).
\] (2.94)

(c) The \(R\)-module, or the \([\chi_{1/2}^{1/2} \oplus \chi_{3/2}^{3/2}]\) module. This module is built from the highest weight state \(|h\rangle_R\) that is not an eigenstate of the \(G_0\) operator. We define the level of a state in this module as \(N = N_0 + 1/16\), where \(N_0\) is the level of the state above \(|h\rangle_R\). Eqn.(2.66) gives \(N \in \mathbb{Z}/2\) + 1/16 (At integer levels \(j = 1/2\) and \(m = 1/2\), or \(j = 3/2\) and \(m = 3/2\), while at half odd integer levels \(j = 3/2\) and \(m = 1/2\), or \(j = 1/2\) and \(m = 3/2\)). The number of states at level \(N\) is given by \(P_{1/2}(N) + P_{3/2}(N) = 2P_{1/2}(N)\) defined in (2.72). We choose the normalization \(_R\langle h | h \rangle_R = 1\).

The Kac determinant formula for the \(R\)-module is given by

\[
\det(M_N^R) = [\alpha (h + 1/48 - c/24)]^{P_{1/2}(N)} \prod_{r,s} \det[(h - h_{r,s}^R)(h - h_{r,s}^R)]^{P_{1/2}(N-rs/4)},
\] (2.95)

where \(s - r = 1 \mod 4\). The zeros \(h_{r,s}^R\) read

\[
h_{r,s}^R = \frac{c}{24} - \frac{1}{48} + \frac{1}{96} \left( (r + s)\sqrt{2 - c} + (r - s)\sqrt{8 - c} \right)^2.
\] (2.96)

We give a few low lying states in the \(R\)-module. At level 1/16 there are \(2P_{1/2}(1/16) = 2\) states

\[
|\phi_1\rangle = |h\rangle_R, \quad |\phi_2\rangle = \tilde{G}_0 |h\rangle_R,
\] (2.97)
and the Kac determinant can be calculated using Eqn.(2.58) and GCR (2.68)

\[ \det(M_{\frac{1}{16}}^R) = \alpha(h + 1/48 - c/24) . \]  

(2.98)

At level \(9/16 = 1/2 + 1/16\) there are 2 states

\[ |\phi_1\rangle = \tilde{G}_{-\frac{1}{2}}|h\rangle_R, \quad |\phi_2\rangle = \tilde{G}_{-\frac{1}{2}}\tilde{G}_0|h\rangle_R . \]  

(2.99)

The Kac determinant at this level is given by

\[ \det(M_{\frac{9}{16}}^R) = \frac{16}{9} \alpha(h + 1/48 - c/24)(h - h_{1,2})(h - h_{2,1}) . \]  

(2.100)

At level \(17/16 = 1 + 1/16\) we have 2 states

\[ |\phi_1\rangle = L_{-1}|h\rangle_R, \quad |\phi_2\rangle = L_{-1}\tilde{G}_0|h\rangle_R , \]  

(2.101)

\[ |\phi_3\rangle = \tilde{G}_{-1}|h\rangle_R, \quad |\phi_4\rangle = \tilde{G}_{-1}\tilde{G}_0|h\rangle_R . \]  

(2.102)

and the Kac determinant

\[ \det(M_{\frac{17}{16}}^R) = \frac{256}{81} (h + 1/48 - c/24)^2(h - h_{1,2})(h - h_{2,1})(h - h_{1,4})(h - h_{4,1}) . \]  

(2.103)

(d) The \(R^{(\pm)}\)-modules, or the \([\chi_{3/2}^3]\) and \([\chi_{1/2}^1]\) modules. \(P_{1/2}(N)\) gives the number of states at level \(N\) in the \(R^{(\pm)}\)-modules built from the highest weight states \(|h; g^{\pm}\rangle:\)

\[ G_0|h; g^{\pm}\rangle = g^{\pm}|h; g^{\pm}\rangle . \]  

(2.104)

With our normalization of \(G_0\) in Eqn.(2.34),

\[ G_0 = \begin{cases} \tilde{G}_0 & \text{if } c \leq 8 , \\ i\tilde{G}_0 & \text{otherwise} . \end{cases} \]  

(2.105)

The eigenvalues \(g^{\pm}\) are given by

\[ g^{\pm} = \frac{\sqrt{8 - c}}{8\sqrt{6}} \pm \sqrt{h + 1/48 - c/24} . \]  

(2.106)

The corresponding new determinant formulae read

\[ \det(M_N^{R^{(\pm)}}) = \alpha^{Q(R)} \prod_{r,s} [\alpha(g^{\pm} - g_{r,s})(g^{\pm} - g_{s,r})]^{P_{1/2}(N-rs/4)} , \]  

(2.107)

where \(r, s\) are the same as in (2.93) and

\[ Q(R) = \sum_n P_{1/2}(N - n^2) . \]  

(2.108)

Here \(n^2 \leq N\).

The zeros of the determinants are given by
\[ g_{r,s} = \frac{\sqrt{8-c}}{8\sqrt{6}} + \frac{1}{\sqrt{96}} \text{sgn}(j-1) \left[ (r+s)\sqrt{2-c} + (r-s)\sqrt{8-c} \right], \tag{2.109} \]

where \(2j = (s-r) \mod 4\).

As an example, we consider the \(R^{(\pm)}\)-modules at the lowest non-trivial level \(N = 9/16\). There is one state in each module,

\[ |\phi^\pm\rangle = \tilde{G}_{\frac{h}{2}} |h; g^\pm\rangle. \tag{2.110} \]

The new determinants can be calculated using Eqn. (2.58) and the GCR (2.68)

\[ \det(M_{R^{(\pm)}}^{9/16}) = \frac{4}{3} \alpha (g^+ - g_{1,2})(g^+ - g_{2,1}). \tag{2.111} \]

The Kac determinant (2.100) is related to the new determinants (2.111) via

\[ \det(M_{R}^{9/16}) = \alpha (h + 1/48 - c/24) \det(M_{R^{(+)}}^{9/16}) \det(M_{R^{(-)}}^{9/16}). \tag{2.112} \]

This factorization generalizes to all levels.

Some of the examples of the Kac and new determinants presented in this section will be useful later for the discussion of the null state structure of the critical \(K = 4\) fractional superstring theory. The Kac spectrum of highest weights \(h_{r,s}\) for the \(K = 4\) case was first given in Ref [8].
III. DERIVATION OF KAC AND NEW DETERMINANT FORMULAE

In this section we derive the Kac and new determinant formulae for an arbitrary (integer) level $K$ FSCA using the BRST operator [9,10] in a particular representation of FSCAs via a non-interacting theory of the $\mathbb{Z}_K$ PF and a single boson with the background charge. Since by definition the Kac and new determinants for a given algebra are representation independent, our derivation is valid for all the representations of FSCAs with modules classified in section II.

In subsection A we deduce the highest weights of the primary fields for which there are null states in the modules. The Kac and new determinants vanish for these values of $h$. In subsection B we derive the conformal dimensions of the null states. The difference between these conformal dimensions and the highest weights of the primary fields gives the lowest levels at which the null states appear in the corresponding modules. There we also derive the multiplicities of the null states at higher levels. This fixes the orders of the zeros of the Kac and new determinants. In subsection C we point out the relation between the Kac and new determinants, and via this relation derive the contributions in the Kac determinants due to the fractional supercurrent zero mode.

A. BRST Operators

We consider the representation of the level $K$ FSCA constructed by turning on the background charge of the $\varphi$ boson in the $SU(2)_K$ WZW model discussed in section II. After the background charge is turned on, the energy-momentum tensor of the $\varphi$ boson reads

$$T_\varphi = -\frac{1}{4}(\partial \varphi)^2 + i\alpha_0 \partial^2 \varphi ,$$

(3.1)

The total energy-momentum tensor of the theory is given by

$$T = T_\varphi + T_{PF} ,$$

(3.2)

where $T_{PF}$ is the energy-momentum tensor of the $\mathbb{Z}_K$ PF. The total central charge is

$$c = c_0 - 24\alpha_0^2 \leq c_0 .$$

(3.3)

Once we turn on the background charge, the $SU(2)_K$ symmetry is broken. However, the off-diagonal currents $J^\pm(z)$ (see Eqn.(2.7)) can be modified so that they remain the spin-1 screening currents:

$$S^\pm(z) = \phi_{\pm 1}^0(z) \cdot \exp(i\alpha_{\pm} \varphi(z)) ,$$

(3.4)

where

$$\alpha_{\pm} = \alpha_0 \pm \left( \alpha_0^2 + \frac{1}{K} \right) .$$

(3.5)
The only remaining symmetry that survives the presence of the background charge is the fractional superconformal symmetry. The fractional supercurrent $G$ commutes with the screening charges $S^{(\pm)}$, i.e.,

$$S^{(\pm)}(z)G(w) \sim \frac{W(w)}{(z-w)^2} + \text{reg.} \ ,$$

(3.6)

where the single pole term is absent ($W(w)$ is some operator) and $G$ is given by

$$G = \frac{\sqrt{K}}{2} \left\{ \varepsilon \partial \varphi - i \alpha_0 (K+2) \partial \epsilon \right\} + \frac{iK(\alpha_+ - \alpha_-)}{K+4} \eta \right\} \ .$$

(3.7)

If we take $\alpha_0$ to zero we will recover Eqn. (2.10).

The currents $T(z)$ and $G(z)$, defined in (3.2) and (3.7), generate the level $K$ FSCA (2.11) for the values of the central charge $c \leq c_0$.

Now we consider the BRST operators [9,10]:

$$Q^{(\pm)}_p \equiv p \prod_{i=1}^p \oint \frac{dz_i}{2\pi i} S^{(\pm)}(z_i) \ ,$$

(3.8)

where the $z_j$ integration contour is inside of the $z_i$ contour for $j > i$, and all contours start and end at $z_1$. Since the screening currents, $S^{(\pm)}(z)$, are dimension one operators and their OPE with $G(z)$ is given by (3.6), we conclude that $Q^{(\pm)}_p$ commute with $T(z)$ and $G(z)$:

$$[Q^{(\pm)}_p,T(z)] = 0 \ , \quad [Q^{(\pm)}_p,G(z)] = 0 \ .$$

(3.9)

From the requirement that the primary fields of Eqn. (2.9) remain primary with respect to $G(z)$ and $T(z)$, we find the following expression for the possible primary fields:

$$\chi^j_{j,n,n'}(z) = \phi^j_{j}(z) \cdot \exp(i\beta^j_{n,n'} \varphi(z)) \ ,$$

(3.10)

where $\beta^j_{n,n'} = \frac{1}{2}(1 - n) \alpha_+ + \frac{1}{2}(1 - n') \alpha_-$ and $n, n'$ are integers satisfying the condition

$$2j = (n' - n) \mod K \ .$$

(3.11)

This condition follows from the requirement that the BRST operators $Q^{(\pm)}_p$ be well-defined on the Fock space (There is no conceptual difference between $Q^{(+)p}$ and $Q^{(-)p}$, so from now on we concentrate only on $Q^{(+)p}$). The operation $Q^{(+)p}\phi^j_j \cdot \exp[i\beta^j_{n,n'} \varphi(0)]$ is well defined if and only if the outer $z_1$ contour closes. By the standard method we have

$$\prod_{i=1}^p \exp[i\alpha_+ \varphi(z_i)] \exp[i\beta^j_{n,n'} \varphi(0)] = \prod_{i<k} (z_i - z_k)^{2\alpha_+} \prod_{i=1}^p z_i^{2\alpha_+} \cdot \exp[i\alpha_+ \sum_{i=1}^p \varphi(z_i) + i\beta^j_{n,n'} \varphi(0)] = ,$$

(3.12)

$$\Psi_1(z_1) \cdots \Psi_1(z_p) \phi^j_j(0) = \prod_{i<k} (z_i - z_k)^{-2/K} \prod_{i=1}^p z_i^{-2j/K} \phi^j_{j+p}(0) + \cdots \ .$$

(3.13)
(The latter equation follows from the fusion rules (2.4) and the conformal dimensions of the PF fields given in (2.4)). Then we make the change of variables \( z_i \rightarrow z_1 \cdot u_i \) for \( i = 2, 3, \ldots, p \) and demand the exponent of \( z_1 \) be an integer (so that the \( z_1 \) contour closes). This gives \( p = n \) and the above condition (3.11).

The conformal weight of the field \( \chi^j_{j,n,n'} \) is given by

\[
\Delta^{(j)}_{n,n'} = \frac{c - c_0}{24} + \frac{1}{96} ( (n + n') \sqrt{c_0 - c} + (n - n') \sqrt{c_1 - c} )^2 + \Delta^j_{j} .
\]

In the next subsection we show that for \( n = r > 0 \) and \( n' = s > 0 \) there are null states in the modules built from the highest weight states \( |\chi^j_{j,r,s}\rangle \). Therefore, Eqn.(3.14) gives the values of the highest weight \( h \) at which the Kac and new determinants vanish. Thus, for \( n = r > 0 \) and \( n' = s > 0 \) Eqn.(3.14) gives Eqn.(2.41).

**B. Conformal Weights and Multiplicities of Null States**

To deduce the conformal weights and multiplicities of the null states appearing in the modules, we consider the BRST mapping of the primary states onto the null states. The following diagram illustrates the BRST mapping:

\[
\begin{array}{c}
\chi^j_{-j,-r,-s} \\
\downarrow \\
\chi^\ell_{-j,-r,-s} \\
\downarrow \\
\chi'^j_{-j,-r,-s} \\
\end{array}
\xleftarrow{Q_r^{(+)}}
\begin{array}{c}
\chi^\ell_{-\ell,r,-s} \\
\downarrow \\
\chi'^\ell_{-\ell,r,-s} \\
\end{array}
\]

The vertical axis measures the conformal dimension of the states: The vertical arrows indicate the action of the \( T \) and \( G \) current modes. The \( T \) current has \( j = m = 0 \) quantum numbers, while \( G \) has \( j = 1 \) and \( m = 0 \). Hence in the vertical direction only the spin \( j \) changes. The BRST operator \( Q_r^{(+)} \) has \( j = 0 \) and \( m = r \). Therefore, in the horizontal direction only the magnetic quantum number \( m \) varies. In particular, the action of \( Q_r^{(+)} \) on a field increases its \( m \) quantum number by \( r \) up to periodicity (2.5). In the diagram, \( \chi^j_{m,n,n'} \equiv \phi^j_m(0) \cdot \exp[i\beta_{n,n'}(0)]|0\rangle \).

Consider the highest weight state \( \chi^\ell_{-\ell,r,-s} \) with \( r, s > 0 \). According to Eqn.(3.11) \( 2\ell = (s + r) \mod K \). The BRST operator \( Q_r^{(+)} \) maps this state onto the state \( \chi^\ell_{-j,-r,-s} \) that is a descendent of the highest weight state \( \chi^j_{-j,-r,-s} \):

\[
\chi^\ell_{-j,-r,-s} = Q_r^{(+)} \chi^\ell_{-\ell,r,-s} .
\]

Here \( -j = (-\ell + r) \mod K = -\frac{(s-r)}{2} \mod (K/2) \) in agreement with the condition \( 2j = (s - r) \mod K \) following from (3.11). Since the operator \( Q_r^{(+)} \) commutes with \( T \) and \( G \), the
positive modes of $T$ and $G$ annihilate $\chi_{j,-r,-s}^\ell$ because they annihilate the primary state $\chi_{-\ell,r,-s}^\ell$. This, in particular, means that $\chi_{j,-r,-s}^\ell$ is a primary and descendent at once, hence it is null. The dimension of this null state is the same as of $\chi_{-\ell,r,-s}^\ell$ and equals $\Delta^{(\ell)}_{r,-s}$.

The modules $[\chi_{j,-r,-s}^\ell]$ and $[\chi_{j,r,s}^\ell]$ (wth $2j = (s - r) \mod K$), built from the highest weight states $\chi_{j,-r,-s}^\ell$ and $\chi_{j,r,s}^\ell$, are dual. The duality means that each state in $[\chi_{j,-r,-s}^\ell]$ has a dual state in $[\chi_{j,r,s}^\ell]$ and structurally these are identical. In particular, the highest weight state $\chi_{j,r,s}^\ell$ is dual to $\chi_{j,-r,-s}^\ell$ in the sense that the two-point correlation $\langle \chi_{-\ell,-r,-s}^\ell | \chi_{j,r,s}^\ell \rangle = 1$: The highest weights of the states $\chi_{j,-r,-s}^\ell$ and $\chi_{j,r,s}^\ell$ are equal. Thus, the null state $\chi_{\ell,-r,-s}^\ell$ appearing in the module $[\chi_{j,-r,-s}^\ell]$ is dual to the null state $\chi_{j,r,s}^\ell$ of the module $[\chi_{j,r,s}^\ell]$. These null states have the same conformal dimensions. The level of the state $\chi_{j,r,s}^\ell$ above the highest weight state $\chi_{j,0}^\ell$, i.e., the lowest level in the module $[\chi_{j,r,s}^\ell]$ at which this null state appears with multiplicity one, is given by

$$N_0^{r,s} = \Delta^{(\ell)}_{r,-s} - \Delta_{j,r,s}^{(j)} = \frac{rs}{K} + \Delta_{\ell} - \Delta_{j}^\ell.$$  

The modified level of this state, that we defined in Eqn.(2.23), is then

$$N_{r,s} = \frac{rs}{K} + \Delta_{\ell}^\ell.$$  

The descendents of the null state $\chi_{j,-r,-s}^\ell$, which is primary, are null states as well. However, those are not primary. There is a one-to-one correspondence between the descendents of the primary state $\chi_{\ell,r,-s}^\ell$ and the descendents of the null state $\chi_{-\ell,r,-s}^\ell$ manifested in the BRST mapping. Thus, the number of the descendents of $\chi_{j,-r,-s}^\ell$ with spin $j'$ at some level $N$ in the module $[\chi_{j,-r,-s}^\ell]$ is the same as the number of the descendents of $\chi_{\ell,r,-s}^\ell$ with the same spin $j'$ at the level $N = N - N_{r,s}$ above the highest weight state $\chi_{-\ell,r,-s}^\ell$ in the module $[\chi_{j,r,s}^\ell]$. The modified level, defined in Eqn.(2.23), corresponding to $\tilde{N}_0$ is then $\tilde{N} = \tilde{N}_0 + \Delta_{\ell} = N - rs/K$. The number of descendents of $\chi_{\ell,r,-s}^\ell$ with spin $j'$ at the level $\tilde{N}$ is given by $P_{\ell}^{j}(\tilde{N}) \equiv P_{\ell}^{j}(\tilde{N})$, where $P_{\ell}^{j}(N)$ are defined in Eqn.(2.23). Because of duality, the counting of descendents of the null state $\chi_{j,r,s}^\ell$ in the module $[\chi_{j,r,s}^\ell]$ is the same as the counting of descendents of the null state $\chi_{-\ell,r,-s}^\ell$ in the module $[\chi_{j,-r,-s}^\ell]$. The orders of the zeros of the Kac and new determinants are precisely given by this counting.

Note that only the highest weight states in the BRST cohomology are the true highest weight states $[\chi_{0}^0]$. In this sense, only the states $\chi_{j,r,s}^\ell$ and $\chi_{j,-r,-s}^\ell$ with $r,s > 0$ are the highest weight states, whereas the state $\chi_{\ell,r,-s}^\ell$ is not. The highest weight states $\chi_{j,r,s}^\ell$ and $\chi_{j,-r,-s}^\ell$ are dual to each other, that is, if $\chi_{j,r,s}^\ell$ is a bra vector in the Fock space, then $\chi_{j,-r,-s}^\ell$ is its corresponding ket vector. The $\varphi$ momenta of these states add up to $2\alpha_0$ to make up for the presence of the background charge so that their two-point correlation function is non-vanishing.

Now we turn to the Kac determinant formulae deriving whose explicit form is our primary goal here. We are ready to write down the Kac determinant formula (up to a positive normalization constant independent of $h$ and $c$) for the module $[\chi_{0}^0]$ of an arbitrary level $K$ FSCA. Since the fractional supercurrent zero mode $G_0$ does not act on the highest weight state in this module, the determinant of inner products of states is a polynomial in $h$ at the
lowest non-trivial level. Then by induction it is a polynomial at all the higher levels. This immediately follows from the counting of the null states derived using the BRST operator. Thus, we arrive at (2.22).

In other modules, however, the zero mode $G_0$ acts on the highest weight states. As pointed out in section II, the determinants of inner products of states are polynomials in $h$ only for some modules, but not for all. In the next subsection we resolve this issue and derive the contribution of the zero mode $G_0$ into the Kac determinant via the relation between the Kac and more fundamental new determinants.

C. Zero Mode Contribution

In this subsection we derive the zero mode contribution. This will complete our derivation of the Kac determinant formulae for the $[\chi_{K/4}]$ and $[\chi_j \oplus \chi_{K/2-j}]$ modules ($0 < j < K/4$) and the new determinant formulae for the $[\chi_j]$ modules ($0 < j \neq K/4$).

Consider the module $[\chi_j]$, $0 < j \neq K/4$. The form of the zero mode algebra is fixed by the OPEs (2.11) and the structure constant $\lambda(c)$ (2.12)

$$\left( G_0^2 - 2a^j \sqrt{c_1 - c} G_0 - (L_0 - \tilde{h}^j) \right) |\chi_j^j\rangle = 0 \, ,$$

where we choose the normalization $N$ of the $G_0$ operator

$$G(z)\chi_j^j(0) = Nz^{-(K+4)/(K+2)}G_0\chi_j^j(0) + ...$$

so that in (3.18) there are only two parameters $a^j$ and $\tilde{h}^j$. The $G_0$ mode is diagonal with respect to the highest weight state $|\chi_j^j\rangle$,

$$G_0|\chi_j^j\rangle = g^j|\chi_j^j\rangle \, ,$$

and Eqn.(3.18) becomes

$$(g^j)^2 - 2a^j g^j \sqrt{c_1 - c} - (h - \tilde{h}^j) = 0 \, ,$$

where $h$ is the highest weight of $|\chi_j^j\rangle$.

To fix the quantities $a^j$ and $\tilde{h}^j$, we consider the eigenvalues $g_{r,s}$ of the $G_0$ operator corresponding to the highest weight states $|\chi_{j,r,s}^j\rangle$ created by the vertex operators (3.10) in the representation (3.7) of FSCA via the $Z_K$ PF and a single boson with background charge. We have

$$G_0|\chi_{j,r,s}^j\rangle = g_{r,s}|\chi_{j,r,s}^j\rangle \, .$$

We determine the form of $g_{r,s}$ from the explicit representation of the $G(z)$ current via (3.7) and the primary fields $\chi_{j,r,s}^j$ via (3.10):

$$g_{r,s} = u\sqrt{c_1 - c} + vb_{r,s} \, .$$

Here we introduced $c$-independent quantities $u$ and $v$, and
\[
\begin{align*}
b_{r,s} &\equiv \frac{1}{96} \left[ (r+s)\sqrt{c_0 - c} + (r-s)\sqrt{c_1 - c} \right]. \\
\end{align*}
\] (3.24)

The quantities (3.23) must satisfy Eqn.(3.21) with \( h = \Delta^{(j)}_{r,s} \), where \( \Delta^{(j)}_{r,s} \) is the conformal dimension of the primary field \( \chi^j_{r,s} \) and is given by Eqn.(3.14). This condition completely determines \( a^j \) and \( \tilde{h}^j \):

\[
u = a^j = \frac{|K - 4j|}{\sqrt{24(K + 4)}}, \quad (3.25)
\]

and

\[
\tilde{h}^j = \frac{c}{24} - c (a^j)^2 = \frac{c}{24} - \frac{(j + 1)(K + 2 - 2j)}{3(K + 4)^2}. \quad (3.26)
\]

For the \([\chi^j_0] \) module, \( 0 < j \neq K/4 \), we have \( v = \text{sgn}(j - K/4) \) and we arrive at Eqn.(2.43).

For the \([\chi^{K/4}_{K/4}] \) module \( v = \pm 1 \) for each pair \( r, s \):

\[
g^{(1,2)}_{r,s} = \pm b_{r,s}. \quad (3.27)
\]

Once we have derived \( a^j \) and \( \tilde{h}^j \), we can solve Eqn.(3.21):

\[
g^j = a^j \sqrt{c_1 - c} + \text{sgn}(j - K/4) \sqrt{h - h^j_0}, \quad (3.28)
\]

where

\[
h^j_0 = \frac{c - c_0}{24} + \Delta^j = \frac{c}{24} - \frac{(K - 4j)^2}{8K(K + 2)}. \quad (3.29)
\]

Note that if \( j = K/4 \), the coefficient \( a^j \) vanishes and

\[
h^{K/4}_0 = \tilde{h}^{K/4} \equiv c/24. \quad (3.30)
\]

Thus, in the module \([\chi^{K/4}_{K/4}] \)

\[
g^\pm = \pm \sqrt{h - c/24}. \quad (3.31)
\]

The determinants of inner products of states, corresponding to the eigenvalues \( g^\pm \), read

\[
\det(M_{(\pm)}^{(K/4)}(N)) = \alpha^{Q(K/4)} \prod_{r,s} ((g^\pm - g_{r,s}^{(1)})(g^\pm - g_{r,s}^{(2)}))^\ell(N-rs/K), \quad (3.32)
\]

where \( r, s, \ell, N \) and \( Q(K/4) \) are the same as in Eqn.(2.30).

Note that \( (g^\pm - g_{r,s}^{(1)})(g^\pm - g_{r,s}^{(2)}) = (g^\pm)^2 - (g_{r,s}^{(1)})^2 = h - h_{r,s} \). This means that the determinants

\[
\det(M_{(\pm)}^{(K/4)}(N)) = \alpha^{Q(K/4)} \prod_{r,s} (h - h_{r,s})^\ell(N-rs/K) \quad (3.33)
\]
are the Kac determinants since they are polynomials in the highest weight $h$. If $G_0$ is not diagonal with respect to the highest weight state, there is the additional zero mode contribution in the Kac determinant. This contribution is given by $(h-c/24)^{P_{K/4}(N)}$. Thus, we arrive at Eqn.(2.32).

The new determinant for the module $[\chi^j]$ ($0 < j \neq K/4$) is a polynomial in the $G_0$ eigenvalue $g^j$ and is given by Eqn.(2.37).

Consider the $[\chi^j \oplus \chi^{K/2-j}]$ module ($0 < j < K/4$) built from the highest weight state

$$|h\rangle = \gamma|h; g^j\rangle + \delta|h; g^{K/2-j}\rangle,$$

where $\gamma\delta \neq 0$ and

$$\langle h; g^j|h; g^{j'}\rangle = \delta_{j,j'}.$$

We label the operators creating the states at level $N$ by $V_i$, $i = 1, ..., P_j(N)$, where $P_j(N)$ is the number of states at this level. Thus, the states in the $[\chi^j]$ module have the form

$$|\phi_i\rangle = V_i|h; g^j\rangle.$$

The new determinant reads

$$\text{det}(\mathcal{M}^{(j)}(N)) = \text{det}(\langle \phi_i | \phi_k \rangle) = \text{det}(\langle h; g^j | Z_{ik} | h; g^{j'} \rangle), \quad Z_{ik} \equiv V_i^\dagger V_k.$$

The Kac determinant for the $[\chi^j \oplus \chi^{K/2-j}]$ module is given by the $(2P_j(N)) \times (2P_j(N))$ matrix

$$\text{det}(\mathcal{M}_0^{(j)}(N)) = \begin{vmatrix} \langle h | Z_{ik} | h \rangle & \langle h | Z_{ik} G_0 | h \rangle \\ \langle h | G_0^\dagger Z_{ik} | h \rangle & \langle h | G_0^\dagger Z_{ik} G_0 | h \rangle \end{vmatrix}$$

$$= [\gamma\delta^2(g^+ - g^-)^{P_j(N)} \text{det}(\mathcal{M}^{(j)}(N))] \text{det}(\mathcal{M}^{(K/2-j)}(N))$$

$$= [2\gamma\delta^2\alpha(h - h_0^j)^{P_j(N)} \text{det}(\mathcal{M}^{(j)}(N))] \text{det}(\mathcal{M}^{(K/2-j)}(N)).$$

This completes the derivation of the Kac and new determinant formulae.

Note, that if $h = h_0^j$, the state $G_0|h\rangle$ becomes null and the Kac determinant vanishes, whereas the new determinants are non-zero.
IV. MINIMAL UNITARY SERIES IN FSCA

In this section we deduce the values of $c$ and $h$ for which FSCAs can have unitary representations. Statistical mechanical systems near the second order phase transition are always expected to be described by an effective unitary field theory with a local order parameter. For the remainder in this section we will thus confine our attention to unitary theories. We analyze unitary representations of FSCAs using the Kac determinants (2.26), (2.32) and (2.39). If for a given module the Kac determinant is negative at any level, it means that there are negative-norm states at that level and the representation is not unitary. If the determinant is greater than or equal zero, further investigation is needed to determine whether or not the representation is unitary. We sketch the non-unitarity proof for FSCAs. Our proof is closely parallel to that for the Virasoro and superconformal minimal series [2,4].

We consider an arbitrary level $K$ FSCA. In the region $c_0 < c < c_1$, all $h_{r,s}$ with $r \neq s$ have non-vanishing imaginary parts, while all $h_{r,r} < 0$. For $c = c_1$ all zeros of the Kac determinant are real and satisfy the following inequality

$$h_{r,s} \leq -j(j+1)(K+4)/K(K+2) \leq 0, \quad 2j = (s-r) \text{ mod } K.$$  

This means that in the region $c_0 < c \leq c_1$, $h > 0$ the Kac determinant is non-vanishing and that all of the Kac matrix eigenvalues are positive. Indeed, as $h \to \infty$, the matrix becomes dominated by its diagonal elements that are strictly positive. On the boundary $c = c_0$ we have

$$h_{r,s} = h_{s,r} = (r-s)^2/4K + \Delta_j^j.$$  

Therefore, for $c = c_0$ and $h > 0$ the Kac determinant vanishes at the points $h = n^2/4K + \Delta_j^j$, $n \in \mathbb{N}$, but does not become negative. Thus, the Kac determinant poses no obstacle in principle to having unitary representations of FSCA in the region $c_0 \leq c \leq c_1$, $h > 0$. The only unitary representation with $h = 0$ is the trivial one with $c = 0$, while there are no unitary representations with $h < 0$.

When $c > c_1$ (2.11) shows that the $G$ current is anti-hermitian and the structure constant $\lambda(c)$ is imaginary. Therefore, all FSCAs with $c > c_1$ are necessarily non-unitary unless $c_{111}$, defined in section III, vanishes. When $K = 1$ or 2 this is exactly the case and there are no analogs of non-unitality proof for $c > 25$ representations of the Virasoro algebra and $c > 27/2$ representations of the superconformal algebra.

In the region $0 < c < c_0$, $h > 0$ the Kac determinant is definitely negative at some level except for the points $(c, h)$ lying on the vanishing curves $h = h_{r,s}(c)$ where the determinant becomes zero. Even on these curves, however, all points, but the ones where they intersect, have ghosts. This discrete set of intersection points, where unitary representations of the FSCA are not excluded, occur at the following values of the central charge

$$c_p = c_0 - \frac{6K}{p(p+K)} = \frac{3K}{K+2} \left(1 - \frac{2(K+2)}{p(p+K)}\right), \quad p = 3, 4, ...$$

($p = 2$ is the trivial theory with $c = 0$). For each $p$ the allowed values of $h$ are given by
\[ h_{r,s} = \frac{[(p + K)r - ps]^2 - K^2}{4Kp(p + K)} + \Delta_j, \]  
(4.4)

where \( 1 \leq r < p, \ 1 \leq s < p + K. \)

Thus, we see that the necessary conditions for unitary highest weight representations of the level \( K \) FSCA are either (for \( K = 1, 2 \) there is no upper bound on the central charge)

\[ \frac{3K}{K + 2} \leq c \leq \frac{3(K + 4)^2}{K(K + 2)}, \quad h \geq 0, \]  
(4.5)

or (4.3) and (4.4). That the latter two conditions are also sufficient, \( i.e., \) that there indeed exist unitary representations of FSCAs for these discrete values of \( c \) and \( h, \) was shown in Ref [10] via the GKO coset space construction [12]. The fractional superconformal unitary minimal model with the central charge \( c_p, \ p \geq 3, \) can be realized by the \( SU(2)_K \otimes SU(2)_p/\overline{SU(2)}_K \otimes SU(2)_p \) coset model. An example of a statistical mechanical system with \( K = 4 \) fractional supersymmetry is the tricritical 3-state Potts model.
V. NO-GHOS T THEOREM AND NULL STATE STRUCTURE OF FRACTIONAL SUPERSTRING

A. No-Ghost Theorem for Subcritical $K = 4$ FSS

In this subsection we prove the no-ghost theorem for the space-time bosonic sector of the subcritical spin-$4/3$ fractional superstring (FSS). We generalize the Brower-Thorn proof of the no-ghost theorem for the bosonic string \[3\] using the Kac determinant formulae for the $K = 4$ FSCA presented in section II. Our discussion closely parallels that of Ref \[3\], so we shall be brief.

The physical state conditions for the space-time bosonic sector of the spin-$4/3$ FSS read

$$L_n|\text{Phys}; q\rangle = G_n^{\pm} |\text{Phys}; q\rangle = 0, \quad L_0|\text{Phys}; q\rangle = v_q|\text{Phys}; q\rangle, \quad n > 0,$$

(5.1)

where $v_q$ is the intercept of the physical states $|\text{Phys}; q\rangle$ with the $Z_3$ charge $q$. Since $|\text{Phys}; +1\rangle = G_0^- |\text{Phys}; -1\rangle$, we have $v_{+1} = v_{-1}$.

Our primary interest in this section is the $c = 5$ representation of the spin-$4/3$ FSS realized by three free bosons and the $SO(2,1)_2$ WZW theory \[14\]. This representation has three dimensional flat Minkowski space-time as its target space, i.e., $SO(2,1)$ global Lorentz symmetry, allowing for the particle interpretation of its scattering amplitudes. The tree level scattering of the physical states does not couple them to the spurious states. In this section we show that the bosonic physical spectrum of this theory is free of negative norm states. This, in particular, means that the tree level scattering is unitary.

The Fock space of this string theory consists of the states

$$|\{\lambda\}, \{\rho\}, \chi\rangle \equiv (\alpha_{-\eta_1})^{\lambda_1} \ldots (\alpha_{-\eta_p})^{\lambda_p} (\epsilon_{-r_1})^{\rho_1} \ldots (\epsilon_{-r_s})^{\rho_s} |\chi; q\rangle,$$

(5.3)

where the ordering is the same as in (2.19). Here $\alpha_{-\eta_i}$, $i = 1, \ldots, p$, are the time-like ($\mu = 0$) creation operators of the world-sheet boson $X^\mu(z)$, $\mu = 0, 1, 2$, while $\epsilon_{-r_k}$, $k = 1, \ldots, s$, are the time-like creation operators of the spin-$1/3$ world-sheet field $\epsilon^\mu(z)$. The state $|\chi\rangle$ is created from the state $\exp[ik \cdot X(0)]|0\rangle$ with the momentum $k^\mu$ by the space-like creation operators of the fields $X^i(z)$ and $\epsilon^i(z)$, $i = 1, 2$. Even though the space-like components of $\epsilon^\mu(z)$ couple to the time-like component, $|\chi\rangle$ is a positive norm state since $SO(3)_2$ is free of ghosts and the space-like components do not change when we rotate $SO(3)_2$ to $SO(2,1)_2$ (the bosonic space-like creation operators certainly do not spoil unitarity). The time-like operators $\alpha_{-\eta}$ and $\epsilon_{-r}$ create negative norm states, i.e., ghosts, in the Fock space of the FSS. The goal of the no-ghost theorem is to prove that there are no ghosts among the physical states that satisfy the above physical state conditions.

The states (5.3) are at the level

$$N_0 = \sum_{i=1}^{p} \lambda_i n_i + \sum_{k=1}^{s} \rho_k r_k$$

(5.4)
above the state $|\chi\rangle$. The metric tensor of these states is given by

$$T_{N_0}(\{\lambda^\prime\}, \{\rho^\prime\}; \{\lambda\}, \{\rho\}) = \delta_{\lambda_1,\lambda_1^\prime} \ldots \delta_{\lambda_p,\lambda_p^\prime} \delta_{\rho_1,\rho_1^\prime} \ldots \delta_{\rho_s,\rho_s^\prime} (-1)^{\nu(p,s)},$$

(5.5)

where

$$\nu(p,s) = \sum_{i=1}^p \lambda_i + \sum_{k=1}^s \rho_k.$$  

(5.6)

The Fock space $\mathcal{F}$ that consists of the states (5.3) can be decomposed into the subspace $\mathcal{R}$ of the primary states $|h; q\rangle$, defined as

$$L_0|h; q\rangle = h|h; q\rangle, \quad L_n|h; q\rangle = G^\pm_{n-(1+q)/3}|h; q\rangle = 0, \quad n > 0,$$

(5.7)

and the subspace of spurious states $\mathcal{S}$, defined as the orthogonal complement of $\mathcal{R}$ in $\mathcal{F}$. Thus, $\mathcal{F} = \mathcal{R} \oplus \mathcal{S}$. The subspace $\mathcal{S}$ is spanned by the states

$$L^\lambda_{-n_1} \ldots L^\lambda_{-n_p} (G^\pm_{r_1})^{\rho_1} \ldots (G^\pm_{r_s})^{\rho_s}|h; q\rangle$$

(5.8)

(where the ordering is the same as in (5.3) and at least one of the numbers $\lambda_i$ or $\rho_k$ is non-zero) provided that their Kac determinant is non-vanishing (the latter condition ensures that these states are linearly independent). The physical states then are the states with $h = v_q = N_0 + \Delta_\chi$, where $\Delta_\chi$ is the $L_0$ eigenvalue of $|\chi; q\rangle$.

Consider the physical states with the intercept $v_q$ and the subspace $\mathcal{S}_q$ of the spurious state space $\mathcal{S}$ that consists of the states with the $L_0$ eigenvalue $v_q$ and the $\mathbb{Z}_3$ charge $q$. According to the Kac determinant formulae for the $K = 4$ FSCA, as $h \to -\infty$, the metric tensor of the states in $\mathcal{S}_q$ at level $N_0$ coincides with (5.7). If we now require the condition

$$\det(\mathcal{M}_N) \neq 0 \quad \text{for all} \quad h < v_q - N_0$$

(5.9)

be satisfied at any level $N_0 > 0$, then the counting of states in $\mathcal{S}_q$ at level $N_0$, given by (5.4), is the same as the counting of states (5.3) with the $L_0$ eigenvalue $v_q$. Moreover, the metric of spurious states in $\mathcal{S}_q$ coincides with (5.3) at any level since the Kac determinant is a polynomial in $h$ and the condition (5.9) is satisfied. This, in particular, means that the number of negative norm states in the Fock space with the $L_0$ eigenvalue $v_q$ is the same as the number of negative norm states in $\mathcal{S}_q$. The latter subspace by definition consists only of spurious states, therefore, there are no ghosts among the physical states. This establishes the no-ghost theorem for the physical states with the intercept $v_q$. In addition to the above condition we should also check the level zero for the absence of ghosts because the $G^\pm_0$ zero modes may create negative-norm physical states. Note that the above analysis can be applied to any representation of the spin-4/3 FSS with only one time-like direction. Eqn. (5.9) is then a sufficient condition for the absence of ghosts in the physical spectrum.

The $\mathbb{Z}_3$ charge $q$ is a quantum number conserved in the tree level scattering processes. We refer to the physical states with $q = \pm 1$ as the $V$-sector, and the physical states with $q = 0$ as the $T$-sector. The lowest lying $V$-sector state is a massless vector particle, whereas its $T$-sector counterpart is a tachyon.

The states that are both spurious and physical are the so-called null states, i.e., they have zero norm. The null states in the $V$-sector come from the $S^{(\pm)}$- and $D^{(\pm)}$-submodules; the null states in the $T$-sector are those in the $S^{(0)}$- and $D^{(0)}$-submodules.
Now we turn to the ghost structure in the $V$-sector. If $c < 8$ the only real zeros of the Kac determinant relevant to the $V$-sector appear in the $S^{(\pm)}$-submodules and are given by
\[ h_{r,r} = \frac{c - 2}{24} (1 - r^2) > \frac{1}{4} - \frac{r^2}{4}, \] (5.10)
where $r$ is an odd integer. On the other hand, $N_0 \geq N_{r,r} = r^2/4 + 1/12$ and $v_{\pm 1} - N_0 < v_{\pm 1} - 1/12 - r^2/4$. Therefore, taking into account the inequality (5.10) we conclude that in the $V$-sector the condition (5.9) is always satisfied for intercepts $v_{\pm 1} < 1/3$ in representations with $c < 8$. However, in this sector the zero modes may spoil unitarity. From (5.2) we find that if $|\text{Phys}; -1\rangle$ is a positive-norm physical state, then $|\text{Phys}; +1\rangle$ has the norm
\[ \langle \text{Phys}; +1 | \text{Phys}; +1 \rangle = (v - c/24) \langle \text{Phys}; -1 | \text{Phys}; -1 \rangle, \] (5.11)
that is non-negative if and only if $v \geq c/24$. Thus, we established the absence of ghosts in the $V$-sector for the following range of $v \equiv v_{\pm 1}$ and $c$:
\[ V - \text{sector} : \quad c/24 \leq v \leq 1/3, \quad c \leq 8. \] (5.12)
The case $v = 1/3$ and/or $c = 8$ was included by continuity.

Next we consider the $T$-sector. If $c < 8$, the only real zeros of the Kac determinant relevant to the $T$-sector appear in the $S^{(0)}$-submodule and are given by (5.10) with $r$ being an even integer. We established the absence of ghosts in the $T$-sector for the following values of $v \equiv v_0$ and $c$:
\[ T - \text{sector} : \quad v \leq (10 - c)/8, \quad c \leq 8. \] (5.13)

For $c < 8$ the condition (5.9) is automatically satisfied in the $R$-sector, because none of the zeros of the Kac determinant are real in this region. The zero mode contribution restricts the intercept to be greater than $c/24 - 1/48$. Thus, we established the absence of ghosts in the $R$-sector for the following values of $v_R$ and $c$:
\[ R - \text{sector} : \quad v_R \geq c/24 - 1/48, \quad c \leq 8. \] (5.14)

### B. Null State Structure of Critical FSS

Any consistent string theory is expected to have extra sets of physical null states at the critical central charge. It is yet unclear if fractional superstrings exist as consistent string theories. Nonetheless, the Kac determinant formulae for FSCAs allow us to examine the null state structure of plausible critical FSS with an arbitrary level $K$ world-sheet fractional supersymmetry. In particular, we recover the well-known results for the bosonic ($K = 1$) and superstring ($K = 2$) theories. We analyze the null state structure of the spin-4/3 (i.e., $K = 4$) FSS at the critical central charge and find extra sets of zero-norm physical states.

First we review the discussion of Ref [9]. The zeros of the Kac determinants for the level $K$ FSCA are given by
\[
hr, s = \frac{c - c_0}{24} + \frac{1}{96} \left( (r + s) \sqrt{c_0 - c} + (r - s) \sqrt{c_1 - c} \right)^2 + \Delta_j^2. \tag{5.15}
\]

Here \( r, s \in \mathbb{N} \); \( c_0 = 3K(K + 2) \) and \( c_1 = 24/K + c_0 \); \( \Delta_j^j = j(K - 2j)/K(K + 2) \), where \( 2j = (s - r) \mod K \). The Kac determinant for the modules with \( j \neq 0 \) also vanishes at \( h_0 = (c - c_0)/24 + \Delta_j^j \) (see section III).

The zero \( h = h_{r, s} \) (i.e., the null state) first appears at the level
\[
N_0^{r, s} = \frac{r s}{K} + \Delta_\ell - \Delta_j^j, \tag{5.16}
\]

where \( 2\ell = (s + r) \mod K \). The intercept of a physical null state is given by
\[
v = hr, s + N_0^{r, s}. \tag{5.17}
\]

Regardless of the central charge, there always is a null state at level \( N_0^{1, 1} \) above the highest weight state \( |h_{1, 1}\rangle = |0\rangle \). For \( K > 1 \) the intercept of this state is
\[
v = \frac{2}{K + 2}; \tag{5.18}
\]

for \( K = 1 \) we have \( v = 1 \). This physical null state belongs to the sector of the FSS theory that contains the massless vector particle (in open FSS) or the graviton (in closed FSS).

The second set of physical null states that belong to the same sector is built from the highest weight state \( |h_{1, K+1}\rangle \). Their intercept must be the same as that of Eqn.(5.18). On the other hand, according to Eqn.(5.17)
\[
v = h_{1, K+1} + N_0^{1, K+1}. \tag{5.19}
\]

This set of constraints has a solution only at the critical central charge \[9\]
\[
c_{\text{critical}} = c_1 + c_0 = \frac{6K}{K + 2} + \frac{24}{K}. \tag{5.20}
\]

For the bosonic string theory \( c_{\text{critical}} = 26 \), \( v = 1 \). For the space-time bosonic sector of the superstring theory we also obtain the well-known result: \( c_{\text{critical}} = 15 \), \( v = 1/2 \). In the fermionic, or Ramond sector of the superstring theory there also are two sets of physical null states. First appears at level zero due to the vanishing of the Kac determinant at \( h = c/24 \). This gives the Ramond sector intercept \( v_R = c/24 = 5/8 \). The second set with the same intercept \( v_R \) appears at the level \( N_0^{1, 2} = 1 \) above the highest weight state \( |h_{1, 2} = -3/8\rangle \) (\( c = 15 \)).

Eqn.(5.20) gives \( c_{\text{critical}} = 10 \) for the \( K = 4 \) FSS. According to Eqn.(5.18) the intercept of the \( V \)-sector physical states is \( v = 1/3 \).

In the \( T \)-sector at \( c = 10 \) there also are two sets of physical null states with the intercept \( v = 1/3 \). The first one occurs at level \( N_0^{1, 3} = 2/3 \) in the \( D^{(0)} \)-submodule built from the highest weight state \( |h_{1, 3} = -1/3\rangle \). However, the second set appears at a rather high level \( N_0^{2, 6} = 3 \) in the \( S^{(0)} \)-submodule. It is unclear if this will pose a problem for unitarity. The no-ghost theorem for the critical \( K = 4 \) FSS is needed to answer this question. In any case, it is likely that the \( T \)-sector (that contains the tachyonic ground state) should be projected out in any consistent string model.
There is another possibility if the $T$- and $V$-sectors can have different intercepts. There are two sets of physical null states in the $T$-sector at $c = 10$ and $v = 0$. The first set appears at level $N_{2}^{2.2} = 1$ in the $S^{(0)}$-submodule built from the highest weight state $|h_{2,2} = -1\rangle$, and the other one occurs at level $N_{0}^{1.7} = 5/3$ in the $D^{(0)}$-submodule built from the highest weight state $|h_{1,7} = -5/3\rangle$. The ground state in this case is no longer tachyonic, but massless.

Extra sets of physical null states also occur in the $R$-sector at $c = 10$ and the intercept $v = 3/8$. The first set appears at level $N_{1}^{1.2} = 1/2$ above the highest weight state $|h_{1,2} = -1/8\rangle$, and the other one occurs at level $N_{0}^{1.4} = 1$ above the highest weight state $|h_{1,4} = -5/8\rangle$.

The above analysis of the physical null state structure of the $K = 4$ FSS indicates that at least the $V$- and $R$-sectors can be expected to be free of ghosts at the critical central charge.

VI. CONCLUSIONS

In this paper we presented and derived the Kac and new determinant formulae for an arbitrary integer level $K$ fractional superconformal algebra. Thus, although complicated, the FSCAs can be studied using tools developed in conformal field theories. Now we know the Kac and new determinant formulae for infinitely many algebras, and only for three of them these determinants can be explicitly calculated. These are the (super)Virasoro and the spin-4/3 parafermion current algebras, for which the (generalized) (anti)commutation relations are known. For the rest of the FSCAs the generalized commutation relations have not yet been written down because of the complications due to the non-abelian braiding properties of these algebras.

Since all of the necessary tools for the rational level $K$ FSCAs have been worked out (namely, the $\mathbb{Z}_{p/q}$ parafermion theory, the rational level $K$ string functions and the BRST operators), it is straightforward to generalize the Kac and new determinants to those algebras.

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