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STUDY OF THE VICINITIES OF SUPERCONFORMAL FIXED POINTS IN TWO-DIMENSIONAL FIELD THEORY

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ԳԵՐԿՈՆՖՈՐՄ ԱՆՇԱՐԺ ԿԵՏԵՐԻ ՇՐՋԱԿԱՅՔԵՐԻ ՈՒՍՈՒՄՆԱՍԻՐՈՒԹՅՈՒՆԸ ԵՐԿՉԱՓ ԴԱՇՏԻ ՏԵՍՈՒԹՅՈՒՆՈՒՄ

Խոտորումների տեսության օգնությամբ կատարված է SM_p ,,նվազագույն գերկոնֆորմ տեսություններին համապատասխանող անշարժ կետերի շրջակայքերի վերանորմախմբային հետազոտություն, $P \gg 1$ դեպքում։ Ըստ $1/P$-ի գլխավոր մոտավորությամբ կառուցված է վերանորմավորման խմբի $SM_p \rightarrow SM_{p-2}$ հետագիծ։

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STUDY OF THE VICINITIES OF SUPERCONFORMAL FIXED POINTS IN TWO-DIMENSIONAL FIELD THEORY

A renormgroup analysis of the vicinities of fixed points corresponding to the "minimal" superconformal theories SM_p with $p \gg 1$ is carried out using the perturbation theory. In the main approximation over $1/p$ the field theory is constructed, corresponding to the renormgroup trajectory $SM_p \to SM_{p-2}$

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ИССЛЕДОВАНИЕ ОКРЕСТНОСТЕЙ СУПЕРКОНФОРМНЫХ НЕПОДВИЖНЫХ ТОЧЕК В ДВУХМЕРНОЙ ТЕОРИИ ПОЛЯ

С помощью теории возмущений проведен ренормгрупповой анализ окрестностей неподвижных точек, соответствующих "минимальным" суперконформным теориям SM_p с $p \gg 1$. В главном приближении по $1/p$ построена теория поля, отвечающая траектории ренормгруппы $SM_p \to SM_{p-2}$.

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The minimal superconformal field theories SM_p , p = = 3, 4,..., are described in [2-4] . We present some facts with the aim to fix notations. The central charge of the Neveu-Schwarz-Ramond algebra for SM_p is given by the formula:

$$\hat{C}_p = 1 - \frac{8}{p(p+2)} \quad , \qquad p = 3, 4, \dots \tag{1}$$

$\hat{C}$ is connected with the usual Virasoro central charge by the relation $\hat{C} = \frac{3}{2} C$. In [4] it is shown that the SM_p models exhaust all two-dimensional superconformal solutions of field theory with $\hat{C} < 1$, satisfying the positivity condition. The space of fields $\mathcal{A}$ which form closed algebra with respect to the operator expansion contains primary Neveu-Schwarz superfields $\Phi_{n,m}(Z,\bar{Z}) = \phi_{n,m} + \theta\Psi_{n,m} + \bar{\theta}\bar{\Psi}_{n,m} + i\theta\bar{\theta}\tilde{\Phi}_{n,m}$ ($n = 1, 2 \dots p-1$; $m = 1, 2, \dots, p+1$; $n+m = 0 (\mathrm{mod}\ 2)$; $(Z,\bar{Z}) \equiv (z,\theta,\bar{z},\bar{\theta})$ are superspace coordinates, where z and $\bar{z}$ are even, and θ and $\bar{\theta}$ are odd coordinates) and Ramond fields $R^{(\alpha)}_{n,m}(z,\bar{z})$ ($n = 1, 2, \dots, p-1$; $m = 1, 2, \dots, p+1$; $m+n = 1 (\mathrm{mod}\ 2)$, $\alpha = \pm$. The dimensions (see [4]) of lowest components of superfields and Ramond fields are:

$$\Delta_{n,m} = \bar{\Delta}_{n,m} = \frac{(mp - n(p+2))^2 - 4}{8p(p+2)} + \frac{1}{32}\left(1 - (-)^{n+m}\right), \qquad (2a)$$

while the dimensions of $\Psi_{n,m}$, $\bar{\Psi}_{n,m}$, $\tilde{\Phi}_{n,m}$ are respectively

$$\left(\Delta_{n,m} + \tfrac{1}{2}, \bar{\Delta}_{n,m}\right); \left(\Delta_{n,m}, \bar{\Delta}_{n,m} + \tfrac{1}{2}\right) u \left(\Delta_{n,m} + \tfrac{1}{2}; \bar{\Delta}_{n,m} + \tfrac{1}{2}\right). \qquad (2b)$$

The structure constants of the operator algebra of SM_p models are calculated in [5]. It is essential for what follows that the superconformal classes $[\tilde{\Phi}_{1,n}]$ and $[R_{1,n}]$ form a subalgebra $A_1 \subset A$.

Consider now a fixed point corresponding to the SM_p theory with $p \gg 1$. From (2a,b) one can see that at $n \ll p$ three series of spinless fields $\tilde{\Phi}_{n,n+2}$, $\tilde{\Phi}_{n+2,n}$ and $\partial_z \partial_{\bar{z}} \Phi_{n,n}$ have dimensions close to 1. Therefore one may expect that the theory perturbed by these operators exhibits nontrivial renorm-group behavior already at small (of $1/p$ order) values of coupling constants, which makes it possible to apply standard perturbation theory. Below we'll study the field theory arising at perturbation of the SM_p model by operator $\tilde{\Phi}_{1,3}$. Its dimensions are:

$$\tilde{\Delta}_{1,3} = \Delta_{1,3} + \frac{1}{2} = 1 - \epsilon \ ; \quad \epsilon = \frac{2}{p+2} \qquad (3)$$

Such consideration is self-consistent, since $\tilde{\Phi}_{1,3}$ is the only field from subalgebra A_1 having dimensions close to 1.

Let $\mathcal{H}^{(p)}(z, \bar{z}, g)$ be the action density of such perturbed theory ($\mathcal{H}^{(p)}(z, \bar{z}, 0)$ describes a fixed point SM_p):

$$\tilde{\Phi}(z, \bar{z}, g) = \frac{\partial}{\partial g} \mathcal{H}^{(p)}(z, \bar{z}, g) \ ; \quad \tilde{\Phi}(z, \bar{z}, 0) = (2\Delta_{1,3})^{-1} \tilde{\Phi}_{1,3}(z, \bar{z}) \qquad (4)$$

Following [7] , we choose g such that

$$G(g) \equiv \langle \widetilde{\Phi}(z,\bar{z},g)\,\widetilde{\Phi}(0,0,g) \rangle \big|_{z\bar{z}=1} = 1 . \tag{5}$$

The trace of the energy-momentum tensor can be written as

$$\Theta = T_\mu^\mu = \beta(g)\widetilde{\Phi} , \tag{6}$$

where $\beta(g)$ is the Gell-Mann-Low function.

The first terms of expansion of β -function in g have the form (see, e.g. [6]):

$$\beta(g) = \epsilon g - \frac{1}{2}(2\pi C) g^2 + \cdots , \tag{7}$$

The value of the structure constant $C = \widetilde{C}_{(1,3)(1,3)(1,3)}(2\Delta_{1,3})^{-3}$ can be derived from the general formulae [5] :

$$C = \frac{2}{\sqrt{3}}(1 + O(\epsilon)). \tag{8}$$

Thus

$$\beta(g) = \epsilon g - \frac{2\pi}{\sqrt{3}} g^2 + \ldots \tag{9}$$

We have written out in (9) only the terms having order ϵ^2 at $g \sim \epsilon$. From (9) it follows that there exists a fixed point

$$2\pi g^* = \sqrt{3}\epsilon(1 + O(\epsilon)). \tag{10}$$

It should be noted that the perturbation by operator $\widetilde{\Phi}_{1,3}$ does not violate global supersymmetry, which is obvious from the identity

$$\int \widetilde{\Phi}_{1,3}(z,\bar{z}) d^2 z = -i \int \Phi_{1,3}(Z,\bar{Z}) d^2 z\, d^2\theta \tag{11}$$

therefore not only conformal but also superconformal symmetry restores at point g^* . Zamolodchikov's theorem on C-function decrease [6] allows one to predict preliminarily that the fixed point (10) corresponds to some model SM_q with $q < p$ (the positivity condition in perturbed theory is not broken). Moreover, due to the additional discrete symmetries available in SM_p models [5] under which the considered perturbation is invariant, there follows that $p - q = 0 \pmod 2$. With the use of formula (see [7])

$$\frac{2}{3}(\hat{C}(g^*) - \hat{C}_p) = -12(2\pi)^2 \int_0^{g^*} \beta(g)\, dg \tag{12}$$

we find:

$$\hat{C}(g^*) = \hat{C}_p - g\epsilon^3 + O(\epsilon^4)\,. \tag{13}$$

Comparison with (1) shows that $\hat{C}_{p-2}$ with required accuracy equals $\hat{C}(g^*)$; hence the fixed point g^* is described by the model SM_{p-2} . The β -function slope at fixed point determines anomalous dimension of field $\tilde{\Phi}(z,\bar{z},g^*)$:

$$\Delta^* = 1 - \frac{d\beta}{dg}\Big|_{g=g^*} = 1 + \epsilon + O(\epsilon^2) \tag{14}$$

hence

$$\tilde{\Phi}(z,\bar{z},g^*) = \tilde{\Phi}^{(p-2)}_{3,1}(z,\bar{z}) \ , \tag{15}$$

where index $p-2$ denotes that the corresponding field refers to the SM_{p-2} model. The constructed field theory (denote it $SM_{p,p-2}$) corresponds to the renormgroup trajectory that connects fixed points SM_p and SM_{p-2} . In the ultraviolet region it approaches SM_p , while the infrared asymptotics of this theory is described by the model SM_{p-2} . The equality (15) shows that $SM_{p,p-2}$ may be considered as SM_{p-2} theory perturbed by operator $\tilde{\Phi}^{(p-2)}_{3,1}$.

In what follows we'll study renormalizations of superfields $\Phi_{n,m}$ in perturbed theory $SM_{p,p-2}$ and establish relations between the fields of "asymptotic" theories SM_p and SM_{p-2} , for which we'll use the method from [7] which needs some reformulation allowing to derive benefit from supersymmetry.

Under infinitesimal superscale transformation $\delta\theta = \frac{1}{2}\varepsilon\theta$; $\delta z = \varepsilon z$, superfield φ is transformed as follows:

$$\delta_\varepsilon \varphi = \varepsilon\left(z\frac{\partial}{\partial z} + \frac{1}{2}\theta\frac{\partial}{\partial\theta} + \hat{\mathcal{D}}\right)\varphi \ , \tag{16}$$

where operator $\hat{\mathcal{D}}$ describes the φ field variation at origin. In particular, for action we have

$$H = \int d^2z\, d^2\theta\, \mathcal{H}(z,\bar{z},\theta,\bar{\theta}) \longrightarrow H + \varepsilon\int d^2z\, d^2\theta\,(\hat{\mathcal{D}} - \tfrac{1}{2})\,\mathcal{H} \equiv$$
$$\equiv H + \varepsilon\int d^2z\, d^2\theta\, \Theta(z,\bar{z},\theta,\bar{\theta}). \tag{17}$$

Suppose our theory depends on n-parametric family of "coupling constants" $g^i = (g^1, \cdots, g^n)$. Denote

$$\Phi_i = \frac{\partial}{\partial g^i}\mathcal{H}(z,\bar{z},\theta,\bar{\theta}). \tag{18}$$

If the theory is renormalizable, then the field $\Theta(z,\bar{z},\theta,\bar{\theta})$ from (17) can be presented in the form:

$$\Theta(z,\bar{z}) = \sum_{i=1}^{n}\beta^i(g)\Phi_i(z,\bar{z}). \tag{19}$$

From above-said one can readily obtain the Callan-Symanzik equation:

$$[\sum_{a=1}^{N}\left(z_a\frac{\partial}{\partial z_a}+\frac{1}{2}\theta_a\frac{\partial}{\partial\theta_a}+\hat{\Gamma}_a(g)\right)-\sum_{i=1}^{n}\beta^i(g)\frac{\partial}{\partial g_i}]\times$$
$$\times<\varphi_1(z_1,\bar{z}_1)\ldots\varphi_N(z_N,\bar{z}_N)>=0, \tag{20}$$

where operator $\hat{\Gamma}=\hat{\mathcal{D}}+\sum_{i=1}^{n}\beta^i\frac{\partial}{\partial g^i}$ is called the anomalous-dimenensions matrix. Consistency of (16), (17), (18) and (19) requires a fulfilment for the following equality:

$$\hat{\Gamma}\Phi_i=\sum_{j=1}^{n}\left(\frac{1}{2}\delta_i^j-\frac{\partial\beta^j}{\partial g^i}\right)\Phi_j\,. \tag{21}$$

Return now to the case of perturbation by operator $\tilde{\Phi}_{1,3}$. When calculating the anomalous-dimensions matrix with the use of perturbation theory one should take into account that only fields with close dimensions mix effectively, so it is easier to examine fields $\Phi_{n,n}(z,\bar{z},g)$ with $n\ll p$, having at $g=0$ dimensions

$$\Delta_{n,n}=\frac{n^2-1}{2P(P+2)}\,. \tag{22}$$

In the considered approximation they do not mix with other fields (this can be seen from operator expansion structure [5]), so

$$\hat{\Gamma}(g)\,\Phi_{n,n}(z,\bar{z},g)=\gamma_{n,n}\,\Phi_{n,n}(z,\bar{z},g). \tag{23}$$

Standard perturbation theory gives

$$\frac{\partial}{\partial g}<\Phi_{n,n}(z_1,\bar{z}_1,g)\,\Phi_{n,n}(z_2,\bar{z}_2,g)>\Big|_{g=0}=-\pi\,\tilde{C}^{(1,3)}_{(n,n)(n,n)}\times$$

$$\times \frac{\Gamma^2(\epsilon)\,\Gamma^2(1-2\epsilon)}{\Gamma^2(2\epsilon)\,\Gamma^2(1-\epsilon)}\left(\hat{z}_{12}\,\hat{\bar{z}}_{12}\right)^{\epsilon-2\Delta_{n,n}} + \left\langle \frac{\partial}{\partial g}\Phi_{n,n}(1)\,\Phi_{n,n}(2)\right\rangle\Big|_{g=0} + \tag{24}$$

$$+ \left\langle \Phi_{n,n}(1)\,\frac{\partial}{\partial g}\Phi_{n,n}(2)\right\rangle\Big|_{g=0} \quad ; \quad \hat{z}_{12} = z_1 - z_2 - \theta_1\theta_2 .$$

Explicit dependence of $\Phi_{n,n}$ on g may be chosen such that the r.h.s. of Eq.(24) vanished at $\hat{z}_{12}\,\hat{\bar{z}}_{12} = 1$ (this corresponds to special choice of coordinates in coupling-constant space). Comparison of (24) with (20) yields:

$$\gamma_{n,n}(g) = \Delta_{n,n} + 2\pi\,\tilde{C}^{1,3}_{(n,n)(n,n)}\,g + \cdots \tag{25}$$

At $\quad g = g^* = (2\pi)^{-1}\sqrt{3}\,\epsilon\,(1 + O(\epsilon))$

$$\gamma_{n,n}(g^*) = \frac{n^2-1}{2P(P+2)} + \frac{n^2-1}{4}\epsilon^3 = \frac{n^2-1}{2(P-2)P} + O(\epsilon^4). \tag{26}$$

In (26) we have taken into account that $\tilde{C}^{(1,3)}_{(n,n)(n,n)} = \frac{n^2-1}{4\sqrt{3}}\epsilon^2 + O(\epsilon^4)$ [5] . Thus,

$$\Phi_{n,n}(z,\bar{z},g^*) = \Phi^{(P-2)}_{n,n}(z,\bar{z}). \tag{27}$$

Consider now $\Phi_{n,n+2}$ and $\Phi_{n,n-2}$ fields. In this case it is necessary to take account of mixing with superfield $\tilde{\Phi}_{n,n} = \frac{i}{2\Delta_{n,n}}\mathcal{D}\bar{\mathcal{D}}\Phi_{n,n}$, where $\mathcal{D} = \frac{\partial}{\partial\theta} + \theta\frac{\partial}{\partial z}$ and $\bar{\mathcal{D}} = \frac{\partial}{\partial\bar{\theta}} + \bar{\theta}\frac{\partial}{\partial\bar{z}}$ are covariant derivatives, and the normalization factor $\frac{i}{2\Delta_{n,n}}$ ensures validity of the equality

$$< \tilde{\Phi}_{n,n}(1)\, \tilde{\Phi}_{n,n}(2) >\Big|_{g=0} = (\hat{z}_{12}\hat{\bar{z}}_{12})^{-2(\Delta_{n,n}+\frac{1}{2})} . \qquad (28)$$

By analogy with the previous case, there take place equalities

$$\frac{\partial}{\partial g} < \tilde{\Phi}_{n,n}(1)\, \tilde{\Phi}_{n,n}(2) >\Big|_{g=0} = -2\pi \cdot \frac{\epsilon}{\Delta^2_{n,n}} \tilde{C}^{(1,3)}_{(n,n)(n,n)} (\hat{z}_{12}\hat{\bar{z}}_{12})^{\epsilon-2\Delta_{n,n}-1} +$$

$$+ < \frac{\partial}{\partial g}\tilde{\Phi}_{n,n}(1)\, \tilde{\Phi}_{n,n}(2) >\Big|_{g=0} + < \tilde{\Phi}_{n,n}(1) \frac{\partial}{\partial g} \Phi_{n,n}(2) >\Big|_{g=0} ,$$

$$\frac{\partial}{\partial g} < \tilde{\Phi}_{n,n}(1)\, \Phi_{n,n\pm2}(2) >\Big|_{g=0} = -\frac{\pi\epsilon(\frac{1}{2}-\epsilon-\Delta_{n,n\pm2})}{\Delta_{n,n}(\frac{1}{2}+\epsilon-\Delta_{n,n\pm2})} C^{(1,3)}_{(n,n)(n,n\pm2)} \times \qquad (29)$$

$$\times (\hat{z}_{12}\hat{\bar{z}}_{12})^{\epsilon-\frac{1}{2}-\Delta_{n,n\pm2}} + < \frac{\partial}{\partial g}\tilde{\Phi}_{n,n}(1)\, \Phi_{n,n\pm2}(2) >\Big|_{g=0} + < \tilde{\Phi}_{n,n}(1) \frac{\partial}{\partial g} \Phi_{n,n\pm2}(2) >\Big|_{g=0} ,$$

$$\frac{\partial}{\partial g} < \Phi_{n,n\pm2}(1)\, \Phi_{n,n\pm2}(2) >\Big|_{g=0} = -\frac{2\pi}{\epsilon} \tilde{C}^{(1,3)}_{(n,n\pm2)(n,n\pm2)} (\hat{z}_{12}\hat{\bar{z}}_{12})^{\epsilon-2\Delta_{n,n\pm2}} +$$

$$+ < \frac{\partial}{\partial g}\Phi_{n,n\pm2}(1)\, \Phi_{n,n\pm2}(2) >\Big|_{g=0} + < \Phi_{n,n\pm2}(1) \frac{\partial}{\partial g} \Phi_{n,n\pm2}(2) >\Big|_{g=0} ,$$

where [5]

$$\frac{\epsilon}{\Delta^2_{n,n}} \tilde{C}^{(1,3)}_{(n,n)(n,n)} = \frac{1}{\sqrt{3}} \frac{4}{n^2-1} \frac{1}{\epsilon} + O(1); \quad \tilde{C}^{(1,3)}_{(n,n\pm2)(n,n\pm2)} = \frac{1}{\sqrt{3}} \frac{n\pm3}{n\pm1} + O(\epsilon); \qquad (30)$$

$$\frac{\epsilon(\frac{1}{2}-\epsilon-\Delta_{n,n\pm2})}{2\Delta_{n,n}(\frac{1}{2}+\epsilon-\Delta_{n,n\pm2})} C^{(1,3)}_{(n,n)(n,n\pm2)} = \frac{1}{\sqrt{3}} \frac{n\mp1}{n\pm1} \sqrt{\frac{n\pm2}{n}} \frac{1}{\epsilon} + O(1) ,$$

where terms $\sim\epsilon(\epsilon\ll 1)$ are neglected.

If we choose

$$\frac{\partial}{\partial g}\Phi_{n,n\pm 2}\Big|_{g=0}=\frac{1}{\epsilon}\frac{\pi}{\sqrt{3}}\frac{n\pm 3}{n\pm 1}\Phi_{n,n\pm 2}+\frac{\pi(n\mp 1)^2}{\epsilon\sqrt{3}\,2(n-1)}\tilde{\Phi}_{n,n}\Big|_{g=0},$$

$$\frac{\partial}{\partial g}\tilde{\Phi}_{n,n}\Big|_{g=0}=\frac{4\pi}{\epsilon\sqrt{3}(n^2-1)}\tilde{\Phi}_{n,n}-\sum_{\sigma=\pm}\frac{\pi(n-\sigma)^2}{\epsilon\sqrt{3}\,2(n+1)}\Phi_{n,n+2\sigma}\Big|_{g=0} \tag{31}$$

then the r.h. sides of Eqs.(29) at $\hat{Z}_{12}\hat{\bar{Z}}_{12}=1$ vanish, and the relation

$$\langle\Phi_\alpha(1)\Phi_\beta(2)\rangle\Big|_{\hat{Z}_{12}\hat{\bar{Z}}_{12}=1}=\delta_{\alpha\beta}+O(g^2);\ \alpha,\beta=1,2,3 \tag{32}$$

is valid, where Φ_1, Φ_2, Φ_3 are $\Phi_{n,n-2}$, $\tilde{\Phi}_{n,n}$, $\Phi_{n,n+2}$, respectively. Using the Callan-Symanzik equation (20) for the two-point correlation function $\langle\Phi_\alpha(1)\Phi_\beta(2)\rangle$ with respect to (29) and (32) one can readily obtain the following expression for the anomalous-dimension matrix at $g=g^*$:

$$\gamma^\alpha_\beta=\frac{1}{2}\delta^\alpha_\beta+\epsilon\begin{pmatrix}\frac{n^2-5}{2(n-1)}; & \frac{n+1}{n-1}\sqrt{\frac{n-2}{n}}; & 0\\ \frac{n+1}{n-1}\sqrt{\frac{n-2}{n}}; & \frac{4}{n^2-1}; & \frac{n-1}{n+1}\sqrt{\frac{n+2}{n}}\\ 0; & \frac{n-1}{n+1}\sqrt{\frac{n+2}{n}}; & \frac{5-n^2}{2(n+1)}\end{pmatrix} \tag{33}$$

This matrix's eigenvalues which determine spectrum of anomalous dimensions of fields $\Phi_{n,n\pm 2}(z,\bar{z},g^*)$; $\tilde{\Phi}_{n,n}(z,\bar{z},g^*)$ are:

$$\tilde{\Delta}_+=\frac{1}{2}+\frac{n+1}{2}\epsilon;\quad \tilde{\Delta}_0=\frac{1}{2}+O(\epsilon^2);\quad \tilde{\Delta}_-=\frac{1}{2}-\frac{n-1}{2}\epsilon. \tag{34}$$

Hence at $g = g^*$ each of these fields is equal to certain linear combination $\Phi^{(p-2)}_{n\pm2,n}$ and $\mathcal{D}\bar{\mathcal{D}}\Phi^{(p-2)}_{n,n}$. Analogously one can examine other $\Phi_{n,m}$ fields (although this is more difficult to do from the algebraic point of view). Note, that the Ramond sector can be studied using word for word the method of Ref. [7] , therefore we have not dwelt on it here.

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Р.Г.ПОГОСЯН

ИССЛЕДОВАНИЕ ОКРЕСТНОСТЕЙ СУПЕРКОНФОРМНЫХ НЕПОДВИЖНЫХ ТОЧЕК В ДВУХМЕРНОЙ ТЕОРИИ ПОЛЯ

(на английском языке, перевод Асланян З.Н.)

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