On the property of Cachazo-Intriligator-Vafa
prepotential at the extremum of the superpotential.

A. Dymarsky* and V. Pestun†

Institute of Theoretical and Experimental Physics
B. Cheremushkinskaya 25, 117259 Moscow, Russia

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Abstract

We consider CIV-DV prepotential \( F \) for \( \mathcal{N} = 1 \) \( SU(n) \) SYM theory at the extremum of the effective superpotential and prove the relation

\[
2F - S_i \frac{\partial F}{\partial S_i} = -u_2 \frac{2\Lambda^{2n}}{n^2 - 1}
\]

1 Introduction

In this note we consider the Cachazo-Intriligator-Vafa prepotential \( [1, 2] \) for \( \mathcal{N} = 1 \) \( SU(n) \) SYM theory, which is equal \( [2] \) to the free energy \( F_m \) for the one matrix holomorphic integral with the polynomial potential, defined as a perturbative expansion around the certain extremum of the action.

\[
e^{-F_m} = \lim_{N \to \infty} \frac{\Lambda^{-N^2}}{Vol(U(N))} \int d\bar{\Phi} e^{-\frac{1}{g^2} \text{tr} W(\Phi_0 + \Phi)}
\]  

(1.1)

Here \( \Phi = \Phi_0 + \Phi \) is a complex matrix, and the integral is an analogue of a contour integral in the one dimensional case, so that the real dimension of the integration domain is \( N^2 \). Naively, one can treat the integral just as an ordinary hermitian matrix
integral in the case of the even order potential, but generally speaking, one should explicitly specify the domain of the integration for the integral to be convergent. While we do not study the problem in its full complexity, we yet adopt the prescription \[2\] of the perturbative expansion around some extremum of the action (it becomes an additional argument of \(F_m\)), and in this case we have not to explicitly specify a contour of integration, since the result depend only on a local structure of the action around the expansion point. So, in a perturbative calculation one can treat the integral just as an ordinary hermitian one matrix model with \(U(N)\) symmetry \[2\].

The action \(W(x)\) is a polynomial of degree \(n + 1\). We define

\[
W'(x) = \prod_{i=1}^{n} (x - \alpha_i) \equiv x^n - \sum_{k=2}^{n} g_k x^{n-k}, \quad \sum_{i=1}^{n} \alpha_i = 0 \quad (1.2)
\]

The critical points of the matrix action are given by specifying \[2, 3, 4\] how \(N\) eigenvalues of the matrix \(\Phi_0\) are distributed around the critical points \(\alpha_i\) of the potential \(W(x)\).

Let \(N_i\) is the number of the eigenvalues of \(\Phi_0\) that are equal to \(\alpha_i\). In the planar (quasiclassical) limit \(g_s \to 0\) the result for the free energy can be represented as an expansion over the genus of fat Feynman diagrams for the matrix model

\[
F_m = \sum_{g=0}^{\infty} g_s^{2g-2} F_g(S_i, \alpha_i) \quad (1.3)
\]

with \(S_i = g_s N_i\). We focus on the planar contribution \(g = 0\) to the free energy \(F\), from which the SW solution for the pure \(\mathcal{N} = 2\) gauge theory can be constructed \[2\].

It has the following structure

\[
F_{pl} = \frac{1}{g_s^2} F(S_i \equiv g_s N_i, \alpha_i), \quad F \equiv F_0 \quad (1.4)
\]

with the following leading terms of \(F\) \[1, 2, 3, 4, 5\]:

\[
F(S_i, \alpha_i) = \sum_i W(\alpha_i) S_i - \sum_i \left( \frac{1}{2} S_i^2 \log \frac{S_i}{\Lambda^2 \Delta_i} - \frac{3}{4} S_i^2 \right) - \sum_{i<j} 2 S_i S_j \log \left( \frac{\alpha_{ij}}{\Lambda} \right) + O(S^3) \quad (1.5)
\]

where

\[
a_{ij} \equiv \alpha_i - \alpha_j \quad (1.6)
\]

\[
\Delta_i = \prod_{j \neq i} a_{ij} = W''(\alpha_i) \quad (1.7)
\]

The first piece of (1.4) is the classical contribution, the second is due to the logarithm of the volume of \(U(N_i)\) group and the gaussian integration, and the third is from the jacobian of \(U(N) \to \prod U(N_i)\).

Then, the following object is believed \[1, 2, 3\] to be a superpotential in the \(\mathcal{N} = 1\) effective low energy theory, that is obtained after integration out the adjoint field \(\Phi\):

\[
W_{eff}(S_i, \alpha_i) = \sum_i \frac{\partial F}{\partial S_i} \quad (1.8)
\]
According to [1, 2], to construct the SW solution [7] for the $\mathcal{N} = 2$ theory, one should consider an extremum point $\langle S_i \rangle$ of $W_{\text{eff}}$

$$\frac{W_{\text{eff}}}{\partial S_i} \bigg|_{\langle S_i \rangle} = \sum_j \frac{\partial^2 F}{\partial S_i \partial S_j} \bigg|_{\langle S_i \rangle} = 0 \quad (1.9)$$

since it is believed that in this point $f_{n-1} = -4\Lambda^{2n}$ [1, 2] and matrix model curve $y^2 = W'^2 + f_{n-1}$ becomes the SW curve [7].

This condition (1.9) could be easily explained from the field theory side [2, 6]. It means decoupling of the common $U(1)$ factor in the effective low-energy theory. Unfortunately, in the matrix integral picture the meaning of this condition is not clear yet, in spite of the rapid development of the issue [22]. We hope that it could be interpreted as a condition for a point, where the perturbative calculation in some "right" way (still not understood rigorously now) coincides with the exact, non-perturbative definition of the matrix integral [22] [8].

In this letter we proposed (instead of discuss) another non-trivial relation that holds in this point * $\langle S \rangle$, involving not only the second derivatives, but the value of $F$ and its first derivatives itself, showing that for the matrix integral the point $\langle S \rangle$ is indeed very special one!

This relation is (here $u_2 \equiv \frac{1}{2} \sum \alpha_i^2 = g_2$ due to $\sum \alpha_i = 0$)

$$V \equiv 2F - \sum S_i \frac{\partial F}{\partial S_i} \bigg|_{\langle S_i \rangle} = \frac{2g_2\Lambda^{2n}}{n^2-1} \quad (1.10)$$

This relation was recently discovered by Matone [9] in the $SU(2)$ case in a slightly different form and was checked by the perturbative expansion. We generalize and prove this relation for an arbitrary unitary gauge group.

Recall, that a similar fact about the $\mathcal{N} = 2$ SW effective prepotential [1]

$$\frac{\partial F_{SW}}{\partial \ln \Lambda^2} = 2F_{SW} - \sum a_i \frac{\partial F_{SW}}{\partial a_i} = -8\pi i\beta_0 u_2 \quad (1.11)$$

has a clear RG interpretation in $\mathcal{N} = 2$ SYM theory [13]. Its simple form can also be explained using superconformal Ward identities [14].

It will be very interesting to find a similar interpretation for the identity (1.10) in terms of matrix model RG [15] [16] and to explain simple and very similar to the r.h.s. of (1.11) form of the r.h.s. of (1.10). Probably, it is possible to find some interpretation of this relation on a field theory side, similarly to the interpretation of the Virasoro constraints in terms of the generalized Konishi anomaly [17, 6].

It should be also remarked that the relation (1.10) looks like a consequence of the exactness of quasiclassical approximation. In fact, dropping all terms except the classical (the first) one in (1.5)

$$F^{\text{cl}} = \sum W(\alpha_i)S_i^{\text{cl}} \quad (1.12)$$

*but not in a general point!  
†which also was discovered by same author [10] and after generalized for other groups and proved in all orders in two ways [11, 12]
and using only classical term for $\langle S \rangle$

\[ S_{cl}^i = \frac{\Lambda^{2n}}{\Delta_i} \]  

(1.13)

one could easily get the result

\[ V_{cl} = \sum W(\alpha_i) \frac{\Lambda^{2n}}{\Delta_i} = \frac{\Lambda^{2n}}{2\pi i} \oint \frac{W(z)}{W'(z)}dz = -\frac{2g_2\Lambda^{2n}}{n^2 - 1} \]  

(1.14)

The relation is not the only one that holds exactly in its classical form at the special point $\langle S \rangle$. The classical result for the superpotential that follows from (1.12)

\[ W_{eff} = \sum W(\alpha_i) \]  

(1.15)

is also exact at the special point $\langle S \rangle$.

The property of the classical approximation to be exact is some consequence of the planar limit in (1.1) and, perhaps, it can enlighten the question of the role of the special point $\langle S \rangle$ from the matrix model point of view.

## 2 The first proof of the relation

In this section we present the proof of the relation (1.10). We consider the planar contribution to the free energy that has a structure

\[ F_{pl} = g_s^{-2}F(g_s N_i; \alpha_i) \]  

(2.16)

Differentiating over $g_s$ we see

\[ \partial_{g_s} F_{pl} = (-2g_s^{-3}F + g_s^{-3}\frac{\partial F}{\partial S_i}S_i) = -g_s^{-3}V \]  

(2.17)

From the definition of the $F_{pl}$ as a free energy we know, that its derivative over a parameter is a vacuum expectation value of the correspondingly coupled operator. Since in the planar limit

\[ F_{pl} = -\log \left( \int d\Phi e^{-g_s^{-1}\text{tr}W(\Phi)} \right) \]  

(2.18)

then

\[ -g_s^{-3}V = \partial_{g_s} F_{pl} = -g_s^{-2}\langle \text{tr} W(\Phi) \rangle V = g_s \langle \text{tr} W(\Phi) \rangle \]  

(2.19)

The expectation values can be calculated with the help of the exact algebraic equation that can be obtained in the planar limit for the resolvent [6]:

\[ R(z) \equiv g_s \text{Tr} \left\langle \frac{1}{z - \Phi} \right\rangle \]  

(2.20)

\[ \text{the leading in the expansion series over } \Lambda \]
The loop equation (or Ward identity for the variation $\delta \Phi = \frac{1}{z-\Phi}$) in the planar limit (when the correlation functions for single trace operators factorize like $\langle A_i B_j \rangle = \langle A_i \rangle \langle B_j \rangle$) is the following

$$R(z)^2 = R(z)W'(z) + \frac{1}{4}f_{n-1}(z)$$

(2.21)

Here $f_{n-1}(z)$ is yet an arbitrary polynomial of the order $n-1$. Its $n$ coefficients specify values of $S_i$ through the relation

$$S_i = \frac{1}{2\pi i} \oint_{A_i} R(z)dz$$

(2.22)

Here and below we use the notations of [18] for cycles on complex plane. From the equation (2.21) we read the solution

$$R(z) = \frac{1}{2} \left( W'(z) - \sqrt{W'(z)^2 + f_{n-1}(z)} \right)$$

Via $R(z)$ one can easily calculate the expectation values of the single trace operators $\text{tr } \Phi^k$ like:

$$g_s \langle \text{tr } \Phi^k \rangle = \frac{1}{2\pi i} \oint_{A} R(z)z^k dz$$

(2.23)

$$A = \sum_{i=1}^{n} A_i, \quad \forall \varphi(z) \quad \frac{1}{2\pi i} \oint_{A} dz \varphi(z) = \text{res}_{z=\infty} \varphi(z)$$

(2.24)

For an arbitrary polynomial we have

$$g_s \langle \text{tr } P(\Phi) \rangle = \frac{1}{2\pi i} \oint_{A} P(z)R(z)dz$$

(2.25)

From [1, 2, 6] we know that at the special point $f_{n-1} = -4\Lambda^2 \Delta_i$ and the leading term for $\langle S \rangle_i$ is indeed given by (1.13)

$$\langle S \rangle_i = -\frac{1}{2\pi i} \oint_{A_i} \frac{1}{2} \sqrt{(x + \alpha_i - 2\Lambda^n/\Delta_i)(x + \alpha_i + 2\Lambda^n/\Delta_i)}\Delta_i + ... = \frac{\Lambda^{2n}}{\Delta_i}$$

(2.26)

Now we are near the desired result

$$g_s \langle \text{tr } W(\Phi) \rangle = -\frac{1}{2\pi i} \oint_{C} W(x)\frac{\sqrt{W'(x)^2 - 4\Lambda^{2n}}}{2} = \text{res}_{z=\infty} \frac{W(z)}{W'(z)} \Lambda^{2n}$$

(2.27)

because

$$\text{res}_{z=\infty} \frac{W(z)}{W'(z)} = \text{res}_{z=\infty} \frac{1}{n+1}x^{n+1} - g_2 \frac{1}{n-1}x^{n-1} + ... = g_2 \left( \frac{1}{n+1} - \frac{1}{n-1} \right) = -g_2 \frac{2}{n^2 - 1}$$

(2.28)

This is the end of the first proof.
3 The second proof of the relation

The second proof of (1.10) checks the definition of $S_D \equiv \frac{\partial F}{\partial S}$ as an integral over some non closed contour $B_i$ and determines the r.h.s. of (1.10) up to an overall constant, which was checked by the perturbative calculations in the introduction. Note, that the coincidence between the perturbative matrix model calculation for $\frac{\partial F}{\partial S}$ [2, 3, 4] and the integral of some differential along $B_i$ cycle [1, 5] is not straightforward, and a lot of work is required to check this identity explicitly even in the first few orders.

Let us consider the partial derivative of (1.10) over $g_k$ considering $S_i = \langle S \rangle_i$ as the functions that depend on $g_k$ and $\Lambda$

$$\frac{\partial V}{\partial g_k} = 2 \frac{\partial F}{\partial g_k} + S^i_D \frac{\partial S^i_D}{\partial g_k} - S^i \frac{\partial S^i}{\partial g_k}$$

(3.29)

at the special point $\langle S \rangle$.

First term could be calculated similarly to the derivative over $g_s$:

$$2 \frac{\partial F}{\partial g_k} = -2 g_s \frac{\langle \text{Tr} \Phi^{n-k+1} \rangle}{n-k+1} = \frac{1}{2\pi i} \int_{A} dz \frac{z^{n-k+1}}{n-k+1} (y - W')$$

$$= \int_{A} \frac{dz}{n-k+1} \left( -\frac{4\Lambda^2 z^{n-k+1}}{2W'} + O\left( \frac{1}{z^{2n}} \right) \right) = -\frac{2\Lambda^2}{(n-1)} \delta_{k,2}$$

(3.30)

The values $S = (g_k, \Lambda)$ and $S_D(g_k, \Lambda)$ defined as [2]

$$S^i = \frac{1}{2\pi i} \oint_{A_i} R(x) dx$$

(3.31)

$$S^i_D = \oint_{B_i} R(x) dx$$

(3.32)

$$R(x) = \frac{1}{2} \left( W'(x) - y(x) \right)$$

(3.33)

Note that the non-compact cycles $B_i$ has a nice property $B_1 - B_n = -b_i$ for $i = 1, n - 1$, where compact cycles $b_i$ and $a_i = A_i$ for $i = 1, n - 1$ form the canonical basis of cycles on the hyperelliptic curve $y^2 = W'/2 - 4\Lambda^{2n}$.

On the special curve $f_{n-1} = -4\Lambda^{2n}$ the differential $ydx$ has no residues. That is why $\sum_i \langle S_i \rangle = 0$, and thus we can change the contour from $B_i$ to $b_i$ ($b_n = 0$) in the definition of $S_D$ (3.31) without changing (3.29).

Now the second and the third terms in (3.29) could be easily calculated with a help of the Riemann bilinear relation [19] yielding

$$S^i_D \frac{\partial S^i}{\partial g_k} - S^i \frac{\partial S^i_D}{\partial g_k} = -\frac{1}{8\pi i} \sum_i \left( \oint_{a_i} ydx \oint_{b_i} \frac{\partial y}{\partial g_k} dx - \oint_{a_i} \frac{\partial y}{\partial g_k} dx \oint_{b_i} ydx \right)$$

$$= -\frac{1}{4} \sum \text{res} \left( \oint_{x_0} y(z)dz \right)$$

(3.34)

(3.35)

It is useful to set $x_0$ to coincide with some root of $y(x)$, then this point will belong to the both branches of the curve, and thus the eventual result is twice larger than the one pole result.
At the infinity (near the pole)
\[ \int_{x_0}^{x} y(z) \, dz = W(x) + \frac{4\Lambda^{2n}}{2(n-1)x^{n-1}} + \text{const} + O\left(\frac{1}{x^n}\right), \quad x \to \infty \tag{3.36} \]
and
\[ \frac{\partial y}{\partial g_k}(x) = -x^{n-k} \left( 1 + \frac{4\Lambda^{2n}}{2(W')^2} + O\left(\frac{1}{x^{4n}}\right) \right), \quad x \to \infty \tag{3.37} \]

Multiplying these two quantities by each other one could simply get
\[ \int_{x_0}^{x} y(z) \, dz \frac{\partial y}{\partial g_k}(x) = -\frac{4\Lambda^{2n}}{2x^{k-1}} \left( \frac{1}{n+1} + \frac{1}{n-1} \right) + Q(x) + O\left(\frac{1}{x^{k+2}}\right), \quad x \to \infty \tag{3.38} \]
where \(Q(x)\) is some polynomial. Thus, substituting (3.38) into (3.34) and adding it to (3.39) we get
\[ \frac{\partial V}{\partial g_k} = -\Lambda^{2n} \left( \frac{1}{n-1} - \frac{1}{n+1} \right) \delta_{k,2} \tag{3.39} \]
This is the end of second proof, since the result (3.39) coincides with the partial derivative over \(g_k\) of (1.10).

4 Conclusion

In this note we proposed a new relation which holds for the CIV prepotential at the special \(N = 2\) point \(\langle S \rangle\). As was mentioned above this relation has the similar form to the renormalization group equation for the Seiberg-Witten prepotential, and it would be interesting to understand it from the matrix model RG point of view.

We expect the straightforward generalization of this relation for other gauge groups \([20]\) and for theory with matter \([21]\).

We hope that the proposed relation will help to specify in some natural way the conditions for the point \(\langle S \rangle\) at the matrix model side, and thus gives more evidence on relation between matrix models and SUSY gauge theories.

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