Three-point correlator of twist-2 operators in BFKL limit

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We compute the correlation function of three twist-2 operators in $\mathcal{N} = 4$ SYM in the leading BFKL approximation at any $N_c$. In this limit, the result is applicable to other gauge theories, including QCD.

INTRODUCTION

The most important physical properties of superconformal $\mathcal{N} = 4$ SYM theory are encoded into its OPE characterized by the spectrum of anomalous dimensions and by the structure constants. While the former is now exactly and efficiently computable at large $N_c$, due to quantum integrability \cite{1}, the calculation of the OPE structure constants is these days on a fast track, especially after the ground-breaking all-loop proposal of \cite{2}.

In this note we calculate the 3-point correlation function of twist-2 operators $O^j(x) = \text{tr} F_{+i} D_{+i}^{-2} F_{+i}^j + \text{fermions} + \text{scalars}$ in $\mathcal{N} = 4$ SYM in the BFKL limit \cite{3} when $\omega = j - 1 \to 0$, the ’t Hooft coupling $g_s \equiv \sqrt{N_c g_{YM}^2} \to 0$ and $g_s^2$ fixed, for arbitrary $N_c$. Since the contribution of fermions+scalars is subleading at this limit, including the internal loops, the result is valid for the pure Yang-Mills theory as well. The case of two-point correlation function was elaborated in our previous paper \cite{4}. In that paper we defined the generalized operators with complex spin as special light-ray operators \cite{5} (regularized as a narrow rectangular Wilson contour called "frame") and calculated their correlation function using OPE over Wilson lines \cite{6} with a rapidity cutoff and the BFKL evolution (see Fig. 1). Here we use the same light-ray operators: one along $n_+$ direction and two along $n_-$. In this case we should use more general Balitsky-Kovchegov (BK) evolution \cite{7, 8} and the leading BFKL contribution comes from the triple-pomeron vertex \cite{9}.

LIGHT-RAY OPERATORS AND THEIR RELATION TO LOCAL OPERATORS

The generalization of local operator $O^j$ for the case of complex spin $j$ was constructed in \cite{4}. It has a form of light-ray operator $\hat{S}^j$ stretched along $n_+$ direction and realizing the principal series representation of $sl(2|4)$ with conformal spin $J = \frac{1}{2} + iv$. The full regularized operator reads as follows:

$$\hat{S}^j(x_\perp) = -(J - 1)(J - 2)\hat{S}_{sc}^j(x_{1\perp}) + \frac{i}{2} \hat{S}_{gl}^j(x_{1\perp}) + \hat{S}_{gl}^j(x_{1\perp}),$$

$$\text{(1)}$$

where for example, the regularized gluon operator is:

$$\hat{S}_{gl}^j(x_{1\perp}) = \lim_{|x_{31\perp}| \to 0} S_{gl}^j(x_{1\perp}, x_{3\perp}) = \lim_{|x_{31\perp}| \to 0} \int_{-\infty}^{\infty} \frac{dx_{1} dx_{3}}{x_{3}^2} \text{tr} F_{+i}^j(x_1) [x_1, x_3]_{\square} F_{+i}(x_3),$$

and $x_1 = (x_{1-}, 0, x_{1\perp})$, $x_3 = (x_{3-}, 0, x_{3\perp})$. Here we introduced the notation $[x_1, x_3]_{\square}$ for rectangular Wilson contour with coordinates $x_1, x_3$ of two diagonally opposite corners, as in Fig. 1. In the case of integer conformal spin $J = j + 1$ it can be rewritten as an integral of local operator $O^j(x)$ with dimension $\Delta(j)$ along a light ray direction $n_+$:

$$\hat{S}^j(x_\perp)|_{x \in \mathbb{N}} = \frac{i^{-j-3/2}}{2^{2j \pi/4}} \int_{-\infty}^{\infty} dx_\perp O^j(x)$$

$$\text{(2)}$$

In this case the correlation function of two light-ray operators stretched along $n_+$ and $n_-$ vectors, normalized as $\langle n_+ n_- \rangle = 1$, is just the double integral of the two-point correlator of local operators $w.r.t.$ the light-ray directions $n_\pm$:

$$\langle \hat{S}^j(x_{1\perp}) \hat{S}^j(y_{1\perp}) \rangle |_{x \in \mathbb{N}} = \frac{b_j}{(|x - y|_{1\perp}^2 \Delta(j) - 1)}$$

$$\text{(3)}$$

FIG. 1. Scheme of computation of 2-point correlator. In the l.h.s., the long sides of regularizing rectangular Wilson frames are stretched along light ray and the short sides in the orthogonal directions. In the r.h.s. we use OPE of frames over color dipoles and compute their correlator, see \cite{4} for details.
The correlator of three light-ray operators is represented, in general, by a sum of several tensor structures. In this note we restrict ourselves to a particular kinematics where all tensor structures have collapsed into a single one. Namely, one light-ray operator is stretched along \( n_+ \) light-ray direction and two other — along \( n_- \). The correlator of 3 light-ray operators can be obtained by integrating the correlation function of 3 local operators along these light-rays. The tensor structures of such local correlators are known from general group-theoretical considerations [10], up to a few structure constants depending on the coupling and symmetry charges. The main dynamical problem which we are addressing here is the calculation of these non-trivial constants. Remarkably, if the coordinates of all 3 light-rays in the space orthogonal to \( \{ n_+, n_- \} \) are restricted to the same line all these structures collapse into a single one, with a single overall structure constant which we are going to compute. It is worth noting that after a conformal transformation the three points in the transverse space take arbitrary positions.

However, the configuration with two collinear light-ray operators is singular, so we first consider three different polarizations \( n_1, n_2, n_3 \) and then take the limit \( n_2 \to n_3 \). The result of integration along light-rays is quite simple and contains only one unknown overall constant

\[
\langle {\hat S}^{n_1} (x_\perp) {\hat S}^{n_2} (y_\perp) {\hat S}^{n_3} (z_\perp) \rangle = C_{\{n_i\}} (\{ \Delta_i \}, \{ j_i \}) = \frac{< n_1 n_2 | j_1, j_2 | n_3 | j_3 >}{(|x-y|^2 \Delta_{1,2,3} (|x-z|^2 \Delta_{1,3,2} (|y-z|^2 \Delta_{2,3,1}))}
\]

(4)

where we used a short-hand notation \( |a|_{j,k} \equiv \frac{1}{2} (a_i + a_j - a_k - 1) \) and \( \{ a_i \} \equiv \{ a_1, a_2, a_3 \} \). In what follows, we assume the existence of a good analytic continuation for \( C_{\{n_i\}} (\{ \Delta j_i \}, \{ j_i \}) \) to non-integer \( j_i \)'s. We take the limit \( n_1 = n_+, n_2 = n_-, n_3 = n_2 \) with the normalization \( n_+ n_- = 1 \). In BFKL regime \( j_1 = 1 + \omega_1 \to 1 \) we obtain:

\[
\langle {\hat S}^{2+\omega_1} (x_\perp) {\hat S}^{2+\omega_2} (y_\perp) {\hat S}^{2+\omega_3} (z_\perp) \rangle = \lim_{n_2 \to n_3} \lim_{n_3 \to n_+} \frac{C_{\{n_i\}} (\{ 1 + \omega_1 \})}{\omega_2 + \omega_3 - \omega_1} \times |x-y|^2 \Delta_{1,2,3} (|x-z|^2 \Delta_{1,3,2} (|y-z|^2 \Delta_{2,3,1})
\]

(5)

where \( \Delta_i = \Delta (1 + \omega_i, g^2) \) is given by BFKL spectrum (see below). We explicitly pulled out the denominator \( \omega_2 + \omega_3 - \omega_1 \geq 0 \) because it will emerge in our forthcoming calculation using the BK evolution. Keeping \( \omega_2 + \omega_3 - \omega_1 \geq 0 \) we 0 \)

\[
\delta (\omega_2 + \omega_3 - \omega_1) \rightarrow \delta (\omega_2 + \omega_3 - \omega_1)
\]

In addition, we keep \( \omega_1 \) positive through the paper.

Finally, the structure constant is normalized using the corresponding 2-point correlators:

\[
C_{\omega_1, \omega_2, \omega_3} = \frac{C_{\{ \Delta_i \}, \{ 1 + \omega_i \}}}{\sqrt{b_1 + \omega_1} b_1 + \omega_2 b_1 + \omega_3}
\]

(6)

**DECOMPOSITION OVER DIPOLES AND BK EVOLUTION**

When calculating the two-point correlator [4] we used a point splitting regularization in orthogonal direction, replacing light-rays by infinitely narrow Wilson frames with inserted fields in the corners (see Fig. 1). Now, for the sake of simplicity, we carry out our calculation for pure Wilson frames, related to our operators of with zero \( R \)-charge in the following way:

\[
\frac{\partial x_{1,1} \partial x_{3,1}}{\partial x_{13,1} \cdot 0, \omega \rightarrow 0} c(g^2 M, N, \omega) S^{2+\omega} (x_1, x_3).
\]

(7)

The coefficient \( c(g^2 M, N, \omega) \) (denoted below as \( c(\omega) \)) depends on the local regularization procedure and at weak coupling it behaves as \( c(\omega) \sim \frac{g^2 \omega}{\omega} \), but its explicit form is irrelevant for us because we are going to calculate the normalized structure constant where it cancels. In general, there are a few types of leading twist-2 operators which appear in this decomposition but in the BFKL limit a single one with the smallest anomalous dimension survives. In addition, in the \( \omega \to 0 \) limit only the term built out of gauge fields alone does contribute [4].

Following the OPE method [6], the pure Wilson frames can be replaced by regularized color dipoles:

\[
[x_1, x_3] \rightarrow N (1 - U^7 (x_{1,1}, x_{3,1})),
\]

(8)

where

\[
U^7 (x_{1,1}, x_{3,1}) = 1 - \frac{1}{N} \text{tr} (U_{x_{1,1} x_{3,1}}^7),
\]

(9)

\[
U_{x_{1,1} x_{3,1}}^7 = P \exp [i g_{YM} \int_{-\infty}^{\infty} dx_+ A_+^\tau (x)],
\]

(10)

\[
A_+^\tau (x) = d^4 k \theta (x_+ - |k_+|) e^{ikx_+} f_\alpha (k),
\]

(11)

where \( f_\alpha \) is a rapidity cutoff. Now we can write:

\[
\langle {\hat S}^{2+\omega_1} (x_{1,1}, x_{3,1}) {\hat S}^{2+\omega_2} (y_{1,1}, y_{3,1}) {\hat S}^{2+\omega_3} (z_{1,1}, z_{3,1}) \rangle = \int_{-\infty}^{\infty} dx_1 \int_{-\infty}^{\infty} dx_3 \int_{-\infty}^{\infty} dx_{31} \times \int dy_1 \int dy_3 \int dy_{31} \int dz_1 \int dz_3 \int dz_{31} \times
\]

(12)

\[
\times \{ U_{y_{1,1} x_{3,1}}^7 \} V_{y_{1,1} y_{3,1}}^\tau V_{y_{1,1} y_{3,1}}^\tau \}
\]

(12)

where \( D_\perp = - \frac{1}{c(\omega_1, c(\omega_2), c(\omega_3))} \partial_{x_{1,1}} \partial_{x_{3,1}} \partial_{y_{1,1}} \partial_{y_{3,1}} \partial_{z_{1,1}} \partial_{z_{3,1}} \).

In our kinematics two dipoles \( V \) and \( W \) have zero \( n_+ \) projection and in the BFKL approximation they form
a "pancake" field configuration in the reference frame related to \( \mathbf{U} \). This means that the rapidity of \( \mathbf{U} \) serves as the upper limit (cutoff) for integrations w.r.t. rapidities of \( \mathbf{V} \) and \( \mathbf{W} \) in our logarithmic approximation. Now one can use the BK evolution to calculate the quantum average in (12). The evolution of the dipole \( \mathbf{U}^{Y_1} \) with respect to rapidity \( Y_1 = e^{\sigma_1} \) can be described by BK equation [7, 8]:

\[
\sigma \frac{d}{d\sigma} \mathbf{U}^{\sigma}(z_1, z_2) = \mathcal{K}_{\text{BK}} \ast \mathbf{U}^{\sigma}(z_1, z_2),
\]

where \( \mathcal{K}_{\text{BK}} \) is an integral operator having the following form in LO approximation:

\[
\mathcal{K}_{\text{LO BK}} \ast \mathbf{U}(z_1, z_2) = \frac{2g^2}{\pi} \int d^2 z_3 \frac{z_2^2}{z_1^2 z_3^2} [\mathbf{U}(z_1, z_3) + \mathbf{U}(z_3, z_2) - \mathbf{U}(z_1, z_2) - \mathbf{U}(z_1, z_3) \mathbf{U}(z_3, z_2)].
\]

Evolution of \( \mathbf{U}^{Y_1} \) goes from \( Y_1 \) to an intermediate \( Y_0 \) w.r.t. the linear part of (13), and then the BK vertex acts at \( Y_0 \) and generates two dipoles which can be contracted with \( \mathbf{V}^{Y_2} \) and \( \mathbf{W}^{Y_3} \). Schematically, it can be written as follows:

\[
\int dY_0 (\mathbf{U}^{Y_1} \to \mathbf{U}^{Y_0}) \otimes \text{(BK vertex at } Y_0) \otimes \left( \langle \mathbf{U}^{Y_0} \mathbf{V}^{Y_2} \rangle \langle \mathbf{U}^{Y_0} \mathbf{W}^{Y_3} \rangle \right)
\]

The linear BFKL evolution of \( \mathbf{U}^{Y_1} \) from \( Y_1 \) to \( Y_0 \) gives:

\[
\mathbf{U}^{Y_1}(x_1, x_3) = \int d\nu \int d^2 x_2 \frac{\nu^2}{\pi^2} \mathcal{E}_{\nu}(x_{10}, x_{30}) e^{\nu(Y_1 - Y_0)},
\]

\[
\cdot \frac{1}{\pi^2} \int d^2 \gamma d^2 \beta \frac{E_{\nu}(\gamma - x_0, \beta - x_0)}{|\gamma - \beta|^4} \mathbf{U}^{Y_0}(\gamma, \beta),
\]

where we denoted \( Y_{ij} = Y_i - Y_j \) and we introduced the function \( \mathcal{E}_{\nu}(z_{10}, z_{20}) = (|z_{10}|^{1/2} |z_{20}|^{1/2})^{1/2 + i\nu} \) which projects dipoles on the eigenstates of BFKL operator with the eigenvalues \( \mathcal{R}(\nu) = 4g^2(2\psi(1) - \psi(1/2 + i\nu) - \psi(1/2 - i\nu)) \). We take here only the sector \( n = 0 \), where \( n \) is the discrete quantum number of \( SL(2, C) \) because the final result depends only on the moduli of distances.

The non-linear part of BK evolution (13) is described by the following renorm group equation:

\[
\frac{\partial}{\partial Y} \mathbf{U}^{Y}(\gamma, \beta) \bigg|_{Y=Y_0} = \frac{2g^2}{\pi} \int d^2 \alpha \frac{\gamma - \beta}{|\gamma - \alpha|^2 |\beta - \alpha|^2} \mathbf{U}^{Y_0}(\gamma, \alpha) \mathbf{U}^{Y_0}(\alpha, \beta)
\]

Finally, we contract the two emerging dipoles \( \mathbf{U}^{Y_0}(\gamma, \alpha) \) and \( \mathbf{U}^{Y_0}(\alpha, \beta) \) with \( \mathbf{V}^{S+}(y_{1\perp}, y_{3\perp}) \) and \( \mathbf{W}^{S+}(z_{1\perp}, z_{3\perp}) \).

Thus for the planar contribution we get:

\[
\langle \mathbf{U}^{Y_1}(x_{1\perp}, x_{3\perp}) \mathbf{V}^{Y_2}(y_{1\perp}, y_{3\perp}) \mathbf{W}^{Y_3}(z_{1\perp}, z_{3\perp}) \rangle_{pl} =
\]

\[
= -\frac{2g^2}{\pi} \int dY_0 \int d\nu_1 \int d^2 x_0 \frac{\nu^2}{\pi^2} \mathcal{E}_{\nu_1}(x_{10}, x_{30}) e^{\nu(Y_1 - Y_0)} \times
\]

\[
\cdot \frac{1}{\pi^2} \int \frac{d^2 \alpha d^2 \beta d^2 \gamma}{|\gamma - \beta|^2 |\gamma - \alpha|^2 |\beta - \alpha|^2} E_{\nu_1}(\gamma - x_0, \beta - x_0).
\]

\[
\cdot \langle \mathbf{U}^{Y_0}(\gamma, \alpha) \mathbf{V}^{Y_2}(y_{1\perp}, y_{3\perp}) \mathbf{W}^{Y_3}(z_{1\perp}, z_{3\perp}) \rangle +
\]

\[
+ \langle \mathbf{U}^{Y_0}(\gamma, \alpha) \mathbf{W}^{Y_3}(z_{1\perp}, z_{3\perp}) \rangle \langle \mathbf{U}^{Y_0}(\alpha, \beta) \mathbf{V}^{Y_2}(y_{1\perp}, y_{3\perp}) \rangle
\]

(18)

The last two terms in (18) give the same contribution so it is enough to know the correlators of two dipoles [4]:

\[
\langle \mathbf{U}^{Y_0}(\gamma, \alpha) \mathbf{V}^{Y_2}(y_{1\perp}, y_{3\perp}) \rangle = -\frac{8g^4(N_c^2 - 1)}{N_c^4} \int d^2 y_0 \cdot
\]

\[
\cdot \int \frac{d\nu_2 \nu_2^2}{(\nu_1^2 + \nu_2^2)^2} E_{\nu_2}(\gamma - y_0, \alpha - y_0) E_{\nu_2}^*(y_{10}, y_{30}) e^{Y_{01}N_c(\nu_2)}
\]

(19)

\[
\langle \mathbf{U}^{Y_0}(\alpha, \beta) \mathbf{W}^{Y_3}(z_{1\perp}, z_{3\perp}) \rangle = -\frac{8g^4(N_c^2 - 1)}{N_c^4} \int d^2 z_0 \cdot
\]

\[
\cdot \int \frac{d\nu_3 \nu_3^2}{(\nu_1^2 + \nu_3^2)^2} E_{\nu_3}(\alpha - z_0, \beta - z_0) E_{\nu_3}^*(z_{10}, z_{30}) e^{Y_{03}N_c(\nu_3)}
\]

(20)

It was argued in [4] that we can make the following identification for rapidities in dipole correlators: \( Y_{02} = \ln \frac{z_{20}}{z_{10}-z_{20}}, Y_{03} = \ln \frac{z_{30}}{z_{10}-z_{30}} \), where \( \lambda \) a cutoff whose precise value is irrelevant in LO. On the other hand, the difference of rapidities of the first dipole and of the BK vertex \( Y_{10} = \ln \frac{z_{10}}{z_{20}} \) corresponds to BFKL evolution. The integral over \( Y_0 = \ln \frac{z_{10}}{\lambda} \) goes from \( Y_1 \) to max(\( Y_2, Y_3 \)). If we plug (18)-(20) into (12) and do the integrals over light ray directions, i.e. over rapidities, we obtain

\[
\langle S^{S+\omega_1}(x_{1\perp}, x_{3\perp}) S^{S+\omega_2}(y_{1\perp}, y_{3\perp}) S^{S+\omega_3}(z_{1\perp}, z_{3\perp}) \rangle_{pl} =
\]

\[
= -\frac{2g^4}{\pi^4 N_c^8} \delta(\omega_1 - \omega_2 - \omega_3) D_{\perp}
\]

\[
\cdot \int d\nu_1 \frac{\nu_1^2}{\pi^2} \frac{1}{\omega_2 + \omega_3 - 8(\nu_1)} \int \frac{d\nu_2 \nu_2^2}{(\nu_1^2 + \nu_2^2)^2} \frac{1}{\omega_2 - 8(\nu_2)}
\]

\[
\cdot \int \frac{d\nu_3 \nu_3^2}{(\nu_1^2 + \nu_3^2)^2} \omega_3 - 8(\nu_3) \int d^2 x_{10} d^2 y_{01} d^2 z_{20} E_{\nu_1}^*(x_{10}, x_{30}) \cdot
\]

\[
\cdot E_{\nu_3}^*(y_{10}, y_{30}) \Upsilon_{pl}(\nu_1, \nu_2, \nu_3; x_{10}, y_{01}, z_{20})
\]

(21)

The delta-function emerges from a similar limit as in the second line of (5) if we replace the vanishing \( \sqrt{\langle n_{21} n_{31} \rangle} \) by the small cutoff in the integration over \( L_1 \). \( \Upsilon_{pl} \) represents
the planar contribution of BK vertex:

\[
\mathcal{Y}_{pl}(\nu_1, \nu_2, \nu_3; x_0, y_0, z_0) = \int \frac{d^2\alpha d^2\beta d^2\gamma}{|\gamma - \beta|^2 |\gamma - \alpha|^2 |\beta - \alpha|^2} E_{\nu_1}(\beta - x_0, \gamma - x_0) \cdot E_{\nu_2}(\alpha - y_0, \gamma - y_0) E_{\nu_3}(\alpha - z_0, \beta - z_0) = \\
\Omega(h_1, h_2, h_3) \left| x_0 - y_0 \right|^4 |h_{1,2,3}| \left| x_0 - z_0 \right|^4 |h_{1,2,3}| \left| y_0 - z_0 \right|^4 |h_{1,2,3}|.
\]

(22)

where \( h_1 = \frac{1}{2} + iv_1, h_2 = \frac{1}{2} + iv_2, h_3 = \frac{1}{2} + iv_3 \) and the function \( \Omega(h_1, h_2, h_3) \) was presented in [11].

Remarkably, we can also take into account the non-planar contribution, which will provide us with the finite \( N_c \) answer for the BKFL structure constant! It appears as a single extra term \( \mathcal{Y}_{npl} \):

\[
\mathcal{Y}_{npl}(\nu_1, \nu_2, h_3; x_0, y_0, z_0) = \int \frac{d^2\beta d^2\gamma}{|\gamma - \beta|^2} E_{\nu_2}(\beta - y_0, \gamma - y_0) E_{\nu_3}(\beta - z_0, \gamma - z_0) = \\
\Lambda(h_1, h_2, h_3) \left| x_0 - y_0 \right|^4 |h_{1,2,3}| \left| x_0 - z_0 \right|^4 |h_{1,2,3}| \left| y_0 - z_0 \right|^4 |h_{1,2,3}|.
\]

(23)

where \( \Lambda(h_1, h_2, h_3) \) was also presented in [11], and the full answer can be obtained from (21) by replacing \( \mathcal{Y}_{pl} \) with \( \mathcal{Y} \) (see in Fig. 2):

\[
\mathcal{Y} = \mathcal{Y}_{pl} - \frac{2 \pi}{N_c^2} \mathcal{Y}_{npl} \text{Re} \left[ \psi(1) + \psi\left(\frac{1}{2} + iv_1\right) - \psi\left(\frac{1}{2} + iv_2\right) - \psi\left(\frac{1}{2} + iv_3\right) \right].
\]

(24)

The integrals over \( x_0, y_0, z_0 \) are easily computable, e.g.

\[
\int d^2x_0 E_{\nu_1}(\beta - x_0, \gamma - x_0) E_{\nu_1}^*(x_0, x_30) = \\
= (\tau^2)^\frac{1}{2} \cdot \left( \frac{1}{2} + iv_1, \frac{1}{2} + iv_2, 1 + 2iv_3, \tau \right) \times \\
\times 2\Gamma(1 + 2iv_1) \pi^2 \Gamma^2\left(\frac{1}{2} + iv_1\right) \Gamma(1 + 2iv_3) \frac{1 + iv_1}{\nu^2} \frac{1 + iv_2}{\nu^2} \frac{1 + iv_3}{\nu^2} G(\nu) + (\nu \leftrightarrow -\nu),
\]

(25)

where \( \tau = \frac{\left|x_1 - x_3\right|^2 \left|\beta - \gamma\right|}{\left|x_1 - \beta\right|^2 \left|x_3 - \gamma\right|^2} \). In the limit \( x_1, x_3 \to x \) we can replace \( \frac{\left|x_1 - x_3\right|^2 \left|\beta - \gamma\right|}{\left|x_1 - \beta\right|^2 \left|x_3 - \gamma\right|^2} \to \frac{\left|x_1 - x_3\right|^2 \left|\beta - \gamma\right|}{\left|x_1 - \beta\right|^2 \left|x_3 - \gamma\right|^2} \to 0 \). For small \( \tau \) we close the \( \nu_1 \) contour in the lower (upper) half-plane for first(second) term, respectively, both of them giving the same contribution. Integrals over \( \alpha, \beta, \gamma \) in (21) can be reduced to \( \mathcal{Y}_{pl} \) which was represented in [11] in terms of Meijer G-functions and hypergeometric functions, and \( \mathcal{Y}_{npl} \) in terms of \( \Gamma \)-functions. Integrals over \( \nu_1 \) can be done by picking up the BFKL poles \( \omega_i = \Omega(\nu_1) \).

Combining (21),(24) and (25) we come to the final expression for 3-point correlation function:

\[
\langle S^{2 + \omega_1}(x_{1,1}, x_{3,1}) S^{2 + \omega_2}(y_{1,1}, y_{1,1}) S^{2 + \omega_3}(z_{1,1}, z_{3,1}) \rangle = \\
= -ig^{10} \delta(\omega_1 - \omega_2 - \omega_3) \int H \Psi(\nu_1, \nu_2, \nu_3) \left| x_1 - y_1 \right|^2 \gamma_1 \gamma_2 \gamma_3 \left| x_1 - x_2 \right|^2 \gamma_1 \gamma_3 \gamma_3 \left| x_1 - y_2 \right|^2 \gamma_1 \gamma_2 \gamma_3,
\]

(27)

where

\[
H = \frac{211(N_c^2 - 1)^2}{\pi N_c^2} \gamma_1(2 + \gamma_1)^3(1 + \gamma_1)(2 + \gamma_2) \times \\
\times (1 + \gamma_2)(2 + \gamma_3)(1 + \gamma_3) G(\nu_1) G(\nu_2) G(\nu_3)
\]

(28)

\[
\gamma_i = \gamma(1 + \omega_i) - \text{anomalous dimension and the coefficient } \Psi(\nu_1, \nu_2, \nu_3) \text{ can be expressed through the functions } \Omega(h_1, h_2, h_3) \text{ and } \Lambda(h_1, h_2, h_3) \text{ defined in (22)-(23) and calculated in [11]}:
\]

\[
\Psi(\nu_1, \nu_2, \nu_3) = \Omega(h_1, h_2, h_3) - \frac{2}{N_c^2} \Lambda(h_1, h_2, h_3).
\]

(29)

where \( h_i^* = \frac{1}{2} + i\nu_i = 1 + \frac{\nu_i}{2} \).

Our final result for normalized structure constant is:

\[
C_{\omega_1, \omega_2, \omega_3} = -ig^2 \frac{\pi}{N_c^2} \left( N_c^2 - 1 \right) \frac{\gamma_1(2 + \gamma_1)^3(1 + \gamma_1)(2 + \gamma_2)}{\pi N_c^2} \Psi(\nu_1, \nu_2, \nu_3).
\]

(30)

Precising the dependence on parameters \( \left\{ \frac{\nu_i}{2} \right\} \), \( g^2 \) and \( N_c \) we can write: \( C_{\omega_1, \omega_2, \omega_3} = \frac{g}{N_c^2} \int \frac{d^2x_1}{\omega_1^2} \frac{d^2x_2}{\omega_2^2} \frac{d^2x_3}{\omega_3^2} \), where \( f \) is a function which depends only on the ratios \( \left\{ \frac{\nu_i}{2} \right\} \).

In the limit \( \frac{\nu_i}{2} \to 0 \) we get the asymptotics:

\[
\Omega(h_1, h_2, h_3) \to \frac{16\pi^3}{\gamma_1^2\gamma_2\gamma_3^3} \gamma_1(2 + \gamma_1)^3(1 + \gamma_1)(2 + \gamma_2) \times \\
\times (1 + \gamma_2)(2 + \gamma_3)(1 + \gamma_3) G(\nu_1) G(\nu_2) G(\nu_3)
\]

(31)

\[
\Lambda(h_1, h_2, h_3) \to \frac{8\pi^2}{\gamma_1^2\gamma_2\gamma_3^3} \gamma_1(2 + \gamma_1)^3(1 + \gamma_1)(2 + \gamma_2) \times \\
\times (1 + \gamma_2)(2 + \gamma_3)(1 + \gamma_3) G(\nu_1) G(\nu_2) G(\nu_3)
\]

(32)

In this limit, the main contribution to 3-point correlator (28) comes from the planar \( O(g^0) \) term

\[
C_{\omega_1, \omega_2, \omega_3} = g^2 \frac{\sqrt{N_c^2 - 1}}{N_c^2} \frac{4i\pi^2}{\omega_1^2\omega_2^2\omega_3^2} (\gamma_2(\omega_2 + \omega_3) + \omega_3(\omega_1 + \omega_3) + \omega_3(\omega_1 + \omega_2) + \omega_1\omega_2\omega_3)(1 + O(g^2))
\]

(33)

whereas the nonplanar one is \( O(g^6) \). It might seem strange that the planar contribution does not start from
It should be possible to generalize our method to the case of multi-point correlators, in particular, for one $n_+$ light-ray operator into any number of $n_-$ operators. We are going to give the answers on some of these questions and clarify many points of our computations in an extended version of this paper [15].

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