EVALUATING THE 6-POINT REMAINDER FUNCTION
NEAR THE COLLINEAR LIMIT

G. PAPATHANASIOU
LAPTh, CNRS, Université de Savoie, Annecy-le-Vieux F-74941, France
and
Physics Department, Theory Division, CERN, CH-1211 Geneva 23, Switzerland

The simplicity of maximally supersymmetric Yang-Mills theory makes it an ideal theoretical
laboratory for developing computational tools, which eventually find their way to QCD ap-
plications. In this contribution, we continue the investigation of a recent proposal by Basso,
Sever and Vieira, for the nonperturbative description of its planar scattering amplitudes, as
an expansion around collinear kinematics. The method of arXiv:1310.5735, for computing
the integrals the latter proposal predicts for the leading term in the expansion of the 6-point
remainder function, is extended to one the subleading terms. In particular, we focus on the
contribution of the 2-gluon bound state in the dual flux tube picture, proving its general form
at any order in the coupling, and providing explicit expressions up to 6 loops. These are
included in the ancillary file accompanying the version of this article on the arXiv.

1 Introduction and Summary

Maximally supersymmetric Yang-Mills theory (MSYM) offers a unique possibility for the non-
perturbative investigation of gauge theories. In its strongly coupled regime it can be mapped to
weakly coupled strings of type IIB on $AdS_5 \times S^5$, which are amenable to perturbative computa-
tions. Furthermore, in the planar limit, where the number of colors $N$ goes to infinity with the
$'t$ Hooft coupling $\lambda \equiv g^2_{YM} N$ fixed, integrable structures emerge, which allow the determina-
tion of certain quantities to all loops $^1$. More importantly, by being the simplest 4-dimensional
interacting gauge theory, it serves as an excellent theoretical laboratory for developing computa-
tional tools, before applying them to QCD. Celebrated examples of this strategy are generalized
unitarity for scattering amplitudes and more recently the method of symbols, for an overview
see $^2$ and references therein. The symbol has been used in calculations of several QCD processes,
such as gluon fusion to heavy quark-antiquark pair $^3$, relevant to experiments at the LHC.

In this contribution, we will focus on the near-collinear kinematics of the planar, Maximally
Helicity Violating (MHV) 6-point amplitude of MSYM. Planarity has the benefit that the only
surviving color structure is a single trace of generators in the adjoint representation of the gauge
group, which we can strip off in order to study its coefficient, the color-ordered amplitude.
Among all different helicity configurations for the external gluons of such amplitudes, it turns
out that the MHV ones $A(+\cdots+\cdots)$, corresponding to all but two helicities being the same,
are the simplest. Remarkably, these amplitudes have been also observed to be dual to Wilson
loops made of straight lightlike segments, as shown in figure 1. And the fact that for $n = 4,5$
legs the dimensionally regulated amplitude is accurately described to all loops by the ansatz
of Anastasiou-Bern-Dixon-Kosower/Bern-Dixon-Smirnov, implies that the 6-point amplitude is
indeed the next interesting case to consider. The remainder function is precisely the part of the
amplitude not captured by the aforementioned ansatz. An account of these developments may be found in \(^2\) as well.

Last but not least, there is growing evidence that each term in the expansion of the remainder function around the limit where consecutive external momenta become collinear, can be computed to all loops with the help of integrability, see \(^4,5,6\) and references therein. For the \(l\)-loop 6-point remainder function \(R_{6}^{(l)}\), symmetry implies that this expansion has the form

\[
R_{6}^{(l)} = \sum_{m=1}^{\infty} e^{-m\tau} \sum_{p=0}^{[m/2]} \cos[(m - 2p)\phi] \sum_{n=0}^{l-1} \tau^n f_{m,p,n}^{(l)}(\sigma)
\]

where \([x]\) denotes the integer part of \(x\), and \(\{\tau, \phi, \sigma\}\) a convenient choice of kinematical variables, in which the collinear limit is described by \(\tau \to \infty\). As we illustrate in figure 1, each term in the sum in \(m\) receives contributions from all \(m\)-particle excitations of a color-electric flux tube, created by the two segments adjacent to the ones becoming collinear, whose dynamics are encoded in an integrable spin chain.

These excitations may also be thought of as insertions the fields of the theory on the side (\(OF\)) of the middle square in figure 1. The leading \(m = 1\) term in (1) comes from a single gluon insertion, and integral expressions for it were found \(^4,5\) by exploiting the aforementioned integrable structures. In \(^1\), we analyzed these integrals, and proved that at any loop order,

\[
f_{1,0,n}^{(l)}(\sigma) = \sum_{s,r,m_i} c_{s,m_1,\ldots,m_r}^{\pm} e^{\pm\sigma s^r} H_{m_1,\ldots,m_r}(-e^{-2\sigma}) , \quad m_i \geq 1 ,
\]

where \(c^{\pm}\) are numeric coefficients, and \(H_{m_1,\ldots,m_n}\) transcendental functions known as harmonic polylogarithms (HPLs). Our proof constituted an algorithm for the direct evaluation of the integrals for arbitrary \(l\), which we employed in order to obtain \(f_{1,0,n}^{(l)}\) for any \(n\) up to \(l = 6\) loops, and for \(n = l - 1\) up to \(n = 12\) loops.

More recently, the \(m = 2\) particle excitations were analyzed, and all-loop integral expressions were also presented for the corresponding term in (1) \(^6\). A variety of different flux tube excitations contributes in this case, and here we will focus on the 2-gluon bound state \(DF\), whose contribution \(W_{DF}\) is part of \(f_{2,0,n}^{(l)}\). Extending the method of \(^7\), we similarly prove that

\[
W_{DF}^{(l)} = \sum_{n=0}^{l-1} \tau^n \tilde{h}_n^{(l)}(\sigma) , \quad \tilde{h}_n^{(l)}(\sigma) = \sum_{s,r,m_i} c_{s,m_1,\ldots,m_r}^{\pm} e^{k\sigma s^r} H_{m_1,\ldots,m_r}(-e^{-2\sigma}) , \quad k = \pm 2, 0 ,
\]

and provide explicit expressions for \(\tilde{h}_n^{(l)}\) up to \(l = 6\) loops. These are included in the computer-readable file \(\text{WDF1-6.m}\) accompanying the version of this article on the \texttt{arXiv}.
The 2-gluon Bound State Contribution

Let us start by reviewing what is known about $R_6$ up to second order in the expansion around collinear kinematics. The kinematical dependence enters through the conformal cross ratios $u_i$ of the cusp positions $x_j$ shown in figure 1, which we parametrize as $(x_0 \equiv x_6)$

$$u_i = \frac{x_{i,j}^2 + x_{i-1,j+3}^2}{x_{i,j}^2 + x_{i+3,j-1}^2}, \quad u_3 = \frac{1}{1 + e^{2\sigma} + 2e^{\sigma-\tau}\cos\phi + e^{-2\tau}}, \quad u_2 = \frac{e^{-\tau}}{2 \cosh \tau}, \quad u_1 = e^{2\sigma + 2\tau} u_2 u_3.$$

Around the collinear limit $\tau \to \infty$, the remainder function has an expansion of the form

$$R_6 = \log W - \log W_{BDS}, \quad W = 1 + e^{-\tau}W_{\text{twist-1}} + e^{-2\tau}W_{\text{twist-2}} + O(e^{-3\tau}),$$

where $W_{BDS}$ is a function known explicitly to all loops, and

$$W_{\text{twist-1}} = \cos \phi W_F, \quad W_{\text{twist-2}} = (W_{\phi\phi} + W_{\psi\bar{\psi}} + W_{FF}) + 2\cos(2\phi)(W_{FF} + W_{DF}),$$

are the contributions of the flux tube excitations of the dual Wilson loop, consisting of gluons $F, \bar{F}$, fermions $\psi, \bar{\psi}$ and scalars $\phi$ of helicity $\pm 1, \pm 1/2$ and 0 respectively. In what follows we will restrict our attention to the contribution of the 2-gluon bound state $DF$,

$$W_{DF} = \int_{-\infty}^{\infty} \frac{du}{2\pi} \mu(u) e^{-\gamma(u)\tau + i p(u)\sigma}.$$

In the last formula, the quantities $\gamma(u), p(u), \mu(u)$ are given to all loops in $g^2 = \lambda/(4\pi)^2$ by

$$\gamma(u) \equiv E(u) - 2 = 4g \sum_{\tau} \cdot M \cdot \kappa(u), \quad p(u) = 2u - 4g \sum_{\tau} \cdot M \cdot \bar{\kappa}(u),$$

$$\mu(u) = \frac{\pi g^2 u (u^2 + 1)}{\sinh(\pi u)(x^{++}x^{--} - g^2)/((x^{++})^2 - g^2)((x^{--})^2 - g^2)} \times \exp \left[ \int_0^\infty \frac{dt}{t} (J_0(2gt) - 1) \frac{2e^{-t} \cos(ut) - J_0(2gt) - 1}{e^t - 1} \right] e^{2\kappa(u) \cdot M \cdot \bar{\kappa}(u) - 2\kappa(u) \cdot Q \cdot M \cdot \kappa(u)}.$$

where $Q$ is a matrix with elements $Q_{ij} = \delta_{ij}(-1)^{i+1}i$, $M$ is related to another matrix $K$,

$$M \equiv (1 + K)^{-1} = \sum_{n=0}^\infty (-K)^n, \quad K_{ij} = 2j(-1)^{j+1} \int_0^\infty \frac{dt}{t} J_i(2gt) J_j(2gt) e^t - 1,$$

$J_i$ is the $i$-th Bessel function of the first kind, and $\kappa, \bar{\kappa}$ are vectors with elements

$$\kappa_j(u) \equiv \int_0^\infty \frac{dt}{t} J_j(2gt) (J_0(2gt) - \cos(ut)) \left[ e^{t/2} \right]^{(-1)^{j-1}} / e^t - 1,$$

$$\bar{\kappa}_j(u) \equiv \int_0^\infty \frac{dt}{t} (-1)^{j+1} J_j(2gt) \sin(ut) \left[ e^{t/2} \right]^{(-1)^{(j+1)-1}} / e^t - 1.$$

Finally, $x^{\pm \pm} = x(u \pm i)$ with $x(u) = (u + \sqrt{u^2 - (2g)^2})/2$.

3 Method and Results

Our main result is the proof that the integral (6) evaluates to the basis (3) at any order $l$ in $g^2 \ll 1$, and the derivation of explicit expressions for $\tilde{h}_n^{(l)}(\sigma)$ up to $l = 6$. To this end, we employ...
the method developed in \(^7\), which consists of reducing the integral into a sum over residues, and using the technology of \(Z\)-sums\(^8\) in order to absorb the summation into the definition of HPLs,

\[
H_{m_1, \ldots, m_r}(x) = \sum_{n_1 > n_2 > \ldots > n_r \geq 1} x^{n_1 \frac{m_1}{n_1} \ldots n_r m_r}.
\]

We have checked that our results for \(\tilde{h}_n^{(l)}(\sigma)\) agree with the expansion of the full \(R_6\) to 4 loops\(^9\), and also with the \(\sigma \to -\infty\) limit given by \(B\) in p.26 of\(^6\). For this we also need to compute \(W_{FF}\) to lowest order, which can be done along similar lines, see also\(^10\). We close by writing a new prediction for part of \(R_6^{(5)}\) all HPLs have argument \(-e^{-2\sigma}\), and \(H_{i,(j,k)} = (H_{i,j,k} + H_{i,k,j})/2\),

\[
\tilde{h}^{(5)}_4 = e^{2\sigma} \left\{ \frac{160}{3} H_{1,(1,3)} + 16 H_{1,2,2} + 32 H_{2,(1,2)} + 32 H_{3,1,1} - 128 H_{1,1,1,1,1} - 2H_5 + (32\sigma - 8) H_{2,1,1,1}\right.
\]
\[
+ \frac{16}{3} H_{1,4} + \frac{32}{3} (H_{2,3} + H_{3,2}) + 64\sigma H_{1,(1,2)} + \left( -32\sigma^2 + 64\sigma - \frac{8\pi^2}{3} - 48 \right) H_{1,1,1}
\]
\[
+ \left( \frac{31}{3} - \frac{32\sigma}{3} \right) H_{3,1} + \left( \frac{40}{3} - \frac{16\sigma}{3} \right) H_{1,3} + \left( \frac{16\sigma}{3} + 8 \right) H_{1,2,2} + \left( 128\sigma - 64 \right) H_{1,1,1,1,1}
\]
\[
+ \left( \frac{1}{2} - 2\sigma \right) H_{1,4} + \left( -\frac{32\sigma^2}{3} + 16\sigma - \frac{8\pi^2}{9} \right) H_{1,2,1} + \left( \frac{32\sigma^2}{3} + 24\sigma + \frac{8}{3} - \frac{8\pi^2}{9} \right) H_{2,1}
\]
\[
+ \left( \frac{2\sigma^2}{3} - \frac{20\sigma}{9} + \frac{4\pi^2}{9} + \frac{4}{3} \right) H_{3} + \left( \frac{4\pi^2}{3} - \frac{2\pi^2}{3} + \frac{28\sigma}{3} + \frac{16\sigma}{3} \right) H_{2,1,1} + \left( \frac{2\sigma^2}{3} + \frac{16\sigma}{9} - \frac{\pi^2\sigma^2}{9} - 12\sigma^2 + \frac{4\pi^2\sigma}{9} + \frac{92\sigma}{3} - \frac{16\sigma\pi(3)}{3} + 8\zeta(3) - \pi^2 - \frac{47}{7} - \frac{7\pi^4}{1050} \right) H_1
\]
\[
+ \left( \frac{32\sigma^3}{3} - 16\sigma^2 + \frac{8\pi^2\sigma}{9} + 48\sigma + 16\zeta(3) - \frac{4\pi^2}{3} - \frac{92}{3} \right) H_{1,1,1} + \frac{40}{3} (H_{1,3} + H_{3,1,1}) + \frac{8}{3} H_4
\]
\[
\left. + \left( -16\sigma^2 + 48\sigma + 4\pi^2 + \frac{92}{3} \right) H_{1,1,1} + 16\sigma (H_{1,2} + H_{2,1}) + (64\sigma - 48) H_{1,1,1,1} + 8H_{2,2,2}
\right.
\]
\[
+ \left( \frac{16\sigma}{3} - 12\sigma + \frac{4\pi^2\sigma}{3} + \frac{92}{3} + 8\zeta(3) - \pi^2 - \frac{47}{3} \right) H_{1} - 64H_{1,1,1,1,1} + \left( \frac{28}{3} - \frac{8\pi^2}{3} \right) H_3
\]
\[
+ \left( \frac{16\sigma^2}{3} + 12\sigma - \frac{4\pi^2}{3} \right) H_2 - \frac{4\pi^2}{3} + \frac{14\sigma^3}{3} + \left( \frac{28}{3} - \frac{47}{3} \right) \sigma^2 + \sigma \left( \frac{-8\zeta(3)}{3} + \frac{7\pi^2}{18} + \frac{47}{3} \right)
\]
\[
+ \frac{20}{3} - \frac{7\pi^2}{9} - \frac{35}{48} - \frac{7\pi^4}{2160} + (\sigma \to -\sigma).
\]

All \(\tilde{h}_n^{(l)}(\sigma)\) up to \(l = 6\) may be found in the ancillary file accompanying this article on the \texttt{arXiv}.

Acknowledgments

We are grateful to B. Basso, J. Drummond and E. Sokatchev for discussions, and to the organizers of the Moriond 2014 for the hospitality and support. This work was supported by the French National Agency for Research (ANR) under contract StrongInt (BLANC-SIMI-4-2011).

References