

# Planar Amplitudes in Maximally Supersymmetric Yang-Mills Theory

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(Dated: September, 2003)

The collinear factorization properties of two-loop scattering amplitudes in dimensionally-regulated  $N = 4$  super-Yang-Mills theory suggest that, in the planar ('t Hooft) limit, higher-loop contributions can be expressed entirely in terms of one-loop amplitudes. We demonstrate this relation explicitly for the two-loop four-point amplitude and, based on the collinear limits, conjecture an analogous relation for  $n$ -point amplitudes. The simplicity of the relation is consistent with intuition based on the AdS/CFT correspondence that the form of the large  $N_c$   $L$ -loop amplitudes should be simple enough to allow a resummation to all orders.

PACS numbers: 11.15.Bt, 11.25.Db, 11.25.Tq, 11.30.Pb, 11.55.Bq, 12.38.Bx

Four-dimensional quantum field theories are extremely intricate, and generically have complicated perturbative expansions in addition to non-perturbative contributions to physical quantities. Gauge theories are interesting in that numerous cancellations occur. This renders perturbative computations more tractable, and their results simpler, than one might otherwise expect. The Maldacena conjecture [1] implies that a special gauge theory is simpler yet: the 't Hooft (planar) limit of maximally supersymmetric four-dimensional gauge theory, or  $N = 4$  super-Yang-Mills theory (MSYM). The conjecture states that the strong coupling limit of this conformal field theory (CFT) is dual to weakly-coupled gravity in five-dimensional anti-de Sitter (AdS) space. The AdS/CFT correspondence is remarkable in taking a seemingly intractable strong coupling problem in gauge theory and relating it to a weakly-coupled gravity theory, which can be evaluated perturbatively. There have been multiple quantitative tests of this correspondence, using observables protected by supersymmetry (see *e.g.* ref. [2]). For unprotected quantities, the domains of validity of the coupling expansions on the gauge and gravity sides do not overlap, so quantitative checks are not possible; however, consistency checks can be established [2].

In this latter context, the AdS/CFT correspondence can be used to motivate a search for patterns in the perturbative expansion of planar MSYM. Intuitively, observables in the strongly coupled limit of this theory should be relatively simple because of the weakly-coupled gravity interpretation. Yet infinite orders in the perturbative expansion, as well as non-perturbative effects, contribute

to the strong coupling limit. How might such a complicated expansion organize itself into a simple result? For quantities protected by supersymmetry, nonrenormalization theorems, or zeros in the perturbative series, are one possibility. Another possibility, for unprotected quantities, is some iterative perturbative structure allowing for a resummation. There have been some hints of an iterative structure developing in the correlation functions of gauge-invariant composite operators [3], but the exact structure, if it exists, is not yet clear.

Amplitudes for scattering of on-shell (massless) quanta — gluons, gluinos, *etc.* — are examples of particular interest because of their importance in QCD applications to collider physics. Although the Maldacena conjecture does not directly refer to on-shell amplitudes, we expect the basic intuition, that the perturbation expansion should have a simple structure, to hold nonetheless. Indeed, the simplicity of one- and two-loop amplitudes in MSYM has allowed their computation to predate corresponding QCD calculations [4, 5].

Perturbative amplitudes in four-dimensional massless gauge theories are not finite, but contain infrared singularities due to soft and collinear virtual momenta. The divergences can be regulated using dimensional regularization with  $D = 4 - 2\epsilon$ . The resulting poles in  $\epsilon$  begin at order  $1/\epsilon^{2L}$  for  $L$  loops, and are described by universal formulæ valid for MSYM, QCD, *etc.* [6]. To preserve supersymmetry we use the four-dimensional helicity scheme [7] variant of dimensional regularization, which is a close relative of dimensional reduction [8]. The infrared divergences turn out to have precisely the iterative struc-

ture we shall find in the full  $N = 4$  amplitudes; thus they provide useful guidance toward exhibiting such a structure.

Infrared divergences generically prevent the definition of a textbook  $S$ -matrix in a non-trivial conformal field theory such as MSYM. For the dimensionally regulated  $S$ -matrix elements we discuss, the regulator explicitly breaks the conformal invariance. However, once the universal infrared singularities are subtracted, the four-dimensional limit of the remaining terms in the amplitudes may be taken, allowing an examination of possible connections to the Maldacena conjecture.

These finite remainders are relevant for computing “infrared-safe” observables in QCD, in which the divergent parts of virtual corrections cancel against real-radiative contributions (not discussed here) to produce finite perturbative results [9]. The finite remainders should also be related to perturbative scattering matrix elements for appropriate coherent states (see *e.g.* ref. [10]). The connection to the  $S$ -matrix for the true asymptotic states of the theory, such as the hadrons of QCD, is of course non-trivial.

In MSYM, there are other hints that higher-loop amplitudes are related in a simple way to the one-loop ones. In particular, the integrands of the amplitudes (prior to evaluation of loop-momentum integrals) have a simple iterative structure [5]. Furthermore, the one-loop amplitudes have a relatively simple analytic structure, which has allowed their determination to an arbitrary number of external legs for configurations with maximal helicity violation [11] and up to six external legs for all helicities [12]. Unitarity then suggests that higher-loop amplitudes may also have a relatively simple analytic structure.

In this Letter we present direct evidence that this intuition is correct for the amplitudes of MSYM. A number of powerful techniques are available to compute them. These include the unitarity-based method [5, 11, 12]; recently-developed multi-loop integration methods (see ref. [13] and references therein); and the imposition of constraints from required behavior as the momenta of two external legs become collinear [14]. Here we shall express the explicit form for the four-point  $N = 4$  amplitude at two loops, in terms of the one-loop amplitude, using previous results [5, 15]. In addition, we present the two-loop splitting amplitude in planar MSYM, computed elsewhere, which summarizes the behavior of amplitudes as the momenta of two legs become collinear. We use the latter to provide evidence that the relationship between the two-loop and one-loop amplitudes continues to hold for an arbitrary number of external legs.

The leading- $N_c$  contributions to the  $L$ -loop  $SU(N_c)$  gauge-theory  $n$ -point amplitudes may be written as,

$$A_n^{(L)} = g^{n-2+2L} N_c^L \sum_{\rho} \text{Tr}(T^{a_{\rho(1)}} T^{a_{\rho(2)}} \dots T^{a_{\rho(n)}})$$

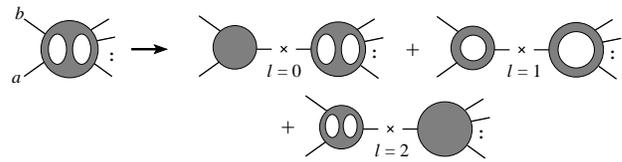


FIG. 1: The collinear factorization of a two-loop amplitude. In each term in the sum, the left blob is a splitting amplitude, and the right blob an  $(n - 1)$ -point scattering amplitude.

$$\times A_n^{(L)}(\rho(1), \rho(2), \dots, \rho(n)), \quad (1)$$

where the sum is over non-cyclic permutations of the external legs, and we have suppressed the momenta and helicities  $k_i$  and  $\lambda_i$ , leaving only the index  $i$  as a label. This decomposition holds for all particles in the gauge super-multiplet because they are all in the adjoint representation.

The color-ordered amplitudes  $A_n^{(L)}(1, 2, \dots, n)$  satisfy simple properties as the momenta of two color-adjacent legs  $k_a, k_b$  become collinear,

$$A_n^{(L)}(\dots, a^{\lambda_a}, b^{\lambda_b}, \dots) \longrightarrow \sum_{l=0}^L \sum_{\lambda=\pm} \text{Split}_{-\lambda}^{(l)}(z; a^{\lambda_a}, b^{\lambda_b}) A_{n-1}^{(L-l)}(\dots, P^\lambda, \dots). \quad (2)$$

The index  $l$  sums over the different loop orders of contributing splitting amplitudes  $\text{Split}_{-\lambda}^{(l)}$ , while  $\lambda$  sums over the helicities of the fused leg  $k_P = -(k_a + k_b)$ , where  $z$  is the momentum fraction of  $k_a$ ,  $k_a = -zk_P$ . The two-loop version of this formula is sketched in fig. 1. The splitting amplitudes are universal and gauge invariant. Formula (2) provides a strong constraint on amplitudes; for example, it has been used to fix the form of a number of one-loop  $n$ -point amplitudes [11, 12, 14].

At tree level, the splitting amplitudes  $\text{Split}_{-\lambda}^{(0)}$  are the same in MSYM as in QCD. Furthermore, the  $N = 4$  supersymmetry Ward identities [16] imply that the MSYM loop splitting amplitudes are all proportional to the tree-level ones, where the ratios depend only on  $z$  and  $\epsilon$ , not on the helicity configuration, nor (except for a trivial dimensional factor) on kinematic invariants [11]. We may therefore write the  $L$ -loop planar splitting amplitudes in terms of “renormalization” factors  $r_S^{(L)}(\epsilon; z, s = (k_1 + k_2)^2)$ , defined by

$$\text{Split}_{-\lambda_P}^{(L)}(1^{\lambda_1}, 2^{\lambda_2}) = r_S^{(L)} \text{Split}_{-\lambda_P}^{(0)}(1^{\lambda_1}, 2^{\lambda_2}). \quad (3)$$

Similarly defining the amplitude ratios  $M_n^{(L)}(\epsilon) \equiv A_n^{(L)}/A_n^{(0)}$ , we obtain in collinear limits,

$$M_n^{(1)}(\epsilon) \rightarrow M_{n-1}^{(1)}(\epsilon) + r_S^{(1)}(\epsilon), \quad (4)$$

$$M_n^{(2)}(\epsilon) \rightarrow M_{n-1}^{(2)}(\epsilon) + r_S^{(1)}(\epsilon)M_{n-1}^{(1)}(\epsilon) + r_S^{(2)}(\epsilon). \quad (5)$$

The  $N = 4$  one-loop splitting amplitudes have been calculated to all orders in  $\epsilon$  [17], with the result

$$r_S^{(1)}(\epsilon; z, s) = \frac{\hat{c}_\Gamma}{\epsilon^2} \left( \frac{\mu^2}{-s} \right)^\epsilon \left[ -\frac{\pi\epsilon}{\sin(\pi\epsilon)} \left( \frac{1-z}{z} \right)^\epsilon + 2 \sum_{k=0}^{\infty} \epsilon^{2k+1} \text{Li}_{2k+1} \left( \frac{-z}{1-z} \right) \right], \quad (6)$$

where  $\text{Li}_n$  is the  $n$ -th polylogarithm,

$$\hat{c}_\Gamma = \frac{e^{\gamma\epsilon}}{2(4\pi)^{2-\epsilon}} \frac{\Gamma(1+\epsilon)\Gamma^2(1-\epsilon)}{\Gamma(1-2\epsilon)}, \quad (7)$$

and  $\gamma$  is Euler's constant. We normalize the amplitudes following the conventions of ref. [6].

We have calculated the two-loop, leading- $N_c$ ,  $N = 4$  splitting amplitudes through  $\mathcal{O}(\epsilon^0)$  using the method of ref. [18] with the result,

$$r_S^{(2)}(\epsilon; z, s) = \frac{1}{2} (r_S^{(1)}(\epsilon; z, s))^2 + f(\epsilon) r_S^{(1)}(2\epsilon; z, s), \quad (8)$$

where

$$f(\epsilon) \equiv (\psi(1-\epsilon) - \psi(1))/\epsilon = -(\zeta_2 + \zeta_3\epsilon + \zeta_4\epsilon^2 + \dots) \quad (9)$$

with  $\psi(x) = (d/dx) \ln \Gamma(x)$ ,  $\psi(1) = -\gamma$ .

The infrared singularities of leading- $N_c$  MSYM at one and two loops can be extracted from more general studies, notably ref. [6]. At one-loop, the divergences are given by,

$$C_n^{(1)}(\epsilon) = -\frac{e^{\gamma\epsilon}}{2\Gamma(1-\epsilon)} \frac{1}{\epsilon^2} \sum_{i=1}^n \left( \frac{\mu^2}{-2k_i \cdot k_{i+1}} \right)^\epsilon. \quad (10)$$

The two-loop divergences, in the four-dimensional helicity scheme, are [6, 19],

$$C_n^{(2)}(\epsilon) = \frac{1}{2} (C_n^{(1)}(\epsilon))^2 + C_n^{(1)}(\epsilon) F_n^{(1)}(\epsilon) - e^{-\epsilon\gamma} \left( \zeta_2 \Gamma(1-2\epsilon) + e^{2\epsilon\gamma} \zeta_3 \epsilon \right) \frac{C_n^{(1)}(2\epsilon)}{\Gamma(1-\epsilon)}. \quad (11)$$

The finite remainder is defined by subtraction,

$$F_n^{(L)}(\epsilon) = M_n^{(L)}(\epsilon) - C_n^{(L)}(\epsilon). \quad (12)$$

Note that  $C_n^{(L)}(\epsilon)$  contains some finite terms as well.

We now present evidence that through  $\mathcal{O}(\epsilon^0)$  the two-loop planar amplitudes are related to one-loop ones via,

$$M_n^{(2)}(\epsilon) = \frac{1}{2} (M_n^{(1)}(\epsilon))^2 + f(\epsilon) M_n^{(1)}(2\epsilon) - \frac{5}{4} \zeta_4. \quad (13)$$

Note the similarity of our ansatz to the two-loop splitting amplitude (8), as well as to the infrared subtraction (11).

The one-loop four-point amplitude in MSYM was first calculated by taking the low energy limit of a superstring [4]. The result is given in terms of a one-loop

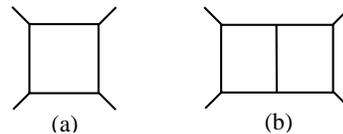


FIG. 2: The scalar integral functions appearing in the (a) one- and (b) two-loop four-point  $N = 4$  amplitudes.

scalar box diagram, depicted in fig. 2(a). (This integral is identical to the one appearing in scalar  $\phi^3$  theory.) Expanding the result in  $\epsilon$  yields,

$$M_4^{(1)}(\epsilon) = \hat{c}_\Gamma \left\{ -\frac{2}{\epsilon^2} \left( \frac{\mu^2}{-s} \right)^\epsilon - \frac{2}{\epsilon^2} \left( \frac{\mu^2}{-t} \right)^\epsilon + \left( \frac{\mu^2}{u} \right)^\epsilon \left[ \frac{1}{2} ((X-Y)^2 + \pi^2) + 2\epsilon \left( \text{Li}_3(x) - X \text{Li}_2(x) - \frac{X^3}{3} - \frac{\pi^2}{2} X \right) - 2\epsilon^2 \left( \text{Li}_4(x) + Y \text{Li}_3(x) - \frac{X^2}{2} \text{Li}_2(x) - \frac{X^4}{8} - \frac{X^3 Y}{6} + \frac{X^2 Y^2}{4} - \frac{\pi^2}{4} X^2 - \frac{\pi^2}{3} X Y - 2\zeta_4 \right) + (s \leftrightarrow t) \right] + \mathcal{O}(\epsilon^3) \right\}, \quad (14)$$

where  $s = (k_1+k_2)^2$ ,  $t = (k_1+k_4)^2$ ,  $u = -s-t$ ,  $x = -s/u$ ,  $y = -t/u$ ,  $X = \ln x$ , and  $Y = \ln y$ . For the four-point case, the  $\epsilon \rightarrow 0$  limit of the finite remainder (12) is

$$F_4^{(1)}(0) = \frac{1}{2} \ln^2 \left( \frac{-s}{-t} \right) + \frac{\pi^2}{2}. \quad (15)$$

In ref. [5] the two-loop  $N = 4$  amplitude was presented in terms of a double-box scalar integral depicted in fig. 2(b), plus its image under the permutation  $s \leftrightarrow t$ . Ref. [15] provides the explicit value of this integral, through  $\mathcal{O}(\epsilon^0)$ , in terms of polylogarithms. Inserting this value, we obtain precisely the result (13) with  $n = 4$ . The equality requires the use of polylogarithmic identities, and involves a non-trivial cancellation of terms between the two contributing integrals. Terms through  $\mathcal{O}(\epsilon^2)$  in  $M_4^{(1)}(s, t)$  contribute at  $\mathcal{O}(\epsilon^0)$  in  $M_4^{(2)}(s, t)$ , since they can multiply the  $1/\epsilon^2$  terms.

Subtracting the two-loop infrared divergence given in eq. (11) from our calculated expression yields

$$F_4^{(2)}(0) = \frac{1}{2} [F_4^{(1)}(0)]^2 - \zeta_2 F_4^{(1)}(0) - \frac{21}{8} \zeta_4, \quad (16)$$

expressed in terms of the one-loop finite remainder (15).

For  $n \geq 5$  legs, we examine the properties as external momenta become collinear, using eq. (2). Applying the one-loop collinear behavior (4) to the ansatz (13), we have

$$M_n^{(2)}(\epsilon) \rightarrow \frac{1}{2} (M_{n-1}^{(1)}(\epsilon) + r_S^{(1)}(\epsilon))^2$$

$$+ f(\epsilon)[M_{n-1}^{(1)}(2\epsilon) + r_S^{(1)}(2\epsilon)] - \frac{5}{4}\zeta_4, \quad (17)$$

which is consistent with the required two-loop collinear properties (5), using eq. (8). Although severely constrained, amplitudes are not uniquely defined by their collinear limits [11]. Thus eq. (13) remains unproven for  $n \geq 5$ . The direct computation of the two-loop five-point function seems feasible, and would provide an important test of the ansatz.

We investigated two potential extensions of the relation (13), each with negative results:

1) We examined the non-planar extension by computing the subleading-color two-loop finite remainders, analogous to  $F_4^{(2)}(0)$ . These terms contain polylogarithms, and hence cannot be written in terms of one-loop finite remainders, unlike the planar eq. (16). Thus the non-planar terms do not appear to have a structure analogous to eq. (13), in line with heuristic expectations from the Maldacena conjecture.

2) For the four-point amplitude, we find that eq. (13) is *not* satisfied at  $\mathcal{O}(\epsilon)$ , due to polylogarithmic obstructions. Hence the relation holds only as  $D \rightarrow 4$ , *i.e.* where the theory becomes conformal.

The possibility of resumming perturbative expansions in MSYM may also have relevance for QCD. QCD may be viewed as containing a ‘‘conformal limit’’ (*e.g.* MSYM) plus conformal-breaking terms. This perspective has had practical impact on topics ranging from the Crewther relation to exclusive processes [20]. We remark that  $N = 4$  amplitudes can be obtained directly from QCD amplitudes by adjusting the number and color of states circulating in the loop: starting from the two-loop QCD amplitudes of ref. [19] and substituting for the ‘spin index dimension’  $D_s = 4 - 2\epsilon\delta_R \rightarrow 10$  and for the color Casimirs  $C_F \rightarrow C_A, T_R N_f \rightarrow 2C_A$ , one obtains the two-loop MSYM amplitudes. (These modifications effectively give  $D = 10, N = 1$  super-Yang-Mills theory, truncated to  $D = 4$ , which is  $N = 4$  super-Yang-Mills theory. See eq. (6.5) of ref. [19] for an analogous conversion to  $D = 4, N = 1$  super-Yang-Mills amplitudes.)

A number of open questions deserve further study. At two loops, the  $N = 4$  planar ansatz should be checked for at least the five-point case. At higher loop order, the intuition described in the introduction suggests a continuation of the iterative structure found at two loops, possibly enabling a resummation of perturbative contributions. Thus we expect that higher-loop planar  $N = 4$  amplitudes will be ‘‘polynomial’’ functions of the one-loop amplitudes. Indeed, the known form of three-loop infrared divergences [6] provide some confirmation of this. Recent advances should make possible explicit evaluation of the leading-color three-loop four-point amplitude, using known expressions for the integrand [5]. (One of the two three-loop integrals needed has already been computed through  $\mathcal{O}(\epsilon^0)$  in terms of generalized polylogarithms [21].) One would also like to identify a symmetry

(presumably related to super-conformal invariance) and associated Ward identity responsible for restricting amplitudes to be iterations of the one-loop amplitude; recall that the relation between the two-loop and one-loop amplitudes holds only near  $D = 4$  where the theory is conformal. We are optimistic that an understanding of the amplitudes of  $N = 4$  super-Yang-Mills theory will lead to new insight into consequences of the AdS/CFT correspondence.

We thank E. D’Hoker, P. Kraus, and E.T. Tomboulis for helpful discussions. This research was supported by the US Department of Energy under contracts DE-FG03-91ER40662 and DE-AC03-76SF00515 and by the *Direction des Sciences de la Matière* of the *Commissariat à l’Energie Atomique* of France.

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