

# A complete solution for scattering in a kind of quiver gauge theory

Su Yingze

Nagoya University

April 21st 2025

# Contents

- 1 Why we need new method?
- 2 Preliminary
- 3 Introduction of quiver gauge theory
- 4 Scattering amplitudes from BCFW

# Why we need new method?

Feynman diagram is a brilliant method without doubt, but it also faces many challenges.

- How to compute high multiplicity amplitudes, like  $n$  gluons scattering.
- Gauge redundancies bring us many structural complications and non-physical degrees of freedom.
- Some symmetries are hidden, like the dual superconformal symmetry for  $\mathcal{N} = 4$  SYM.
- Is the local field theory the most basic description for physics?



# From Frynman diagram to On-shell method

On-shell here means that all quantities we use are gauge invariant and satisfy the on-shell condition. Specifically, there are many ingredients under this frame

- The analytic continuation for S-matrix.
- The color-ordered amplitudes.
- The BCFW recursion relation.
- The spinor helicity discription for amplitudes.

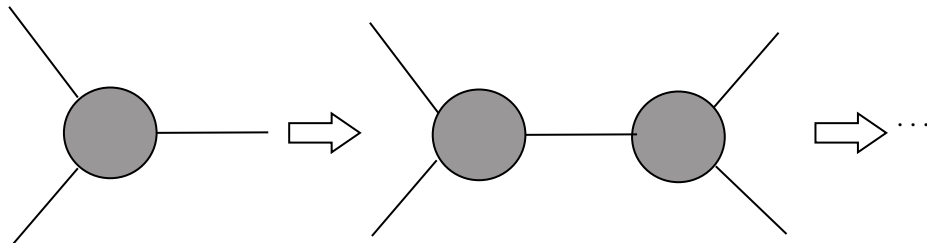
# Contents

- 1 Why we need new method?
- 2 Preliminary**
- 3 Introduction of quiver gauge theory
- 4 Scattering amplitudes from BCFW

# A brief introduction to BCFW

BCFW recursion relation is a method to compute scattering amplitude recursively, means that you can compute higher point amplitudes from lower point.

- Ruth Britto
- Freddy Cachazo
- Bo Feng
- Edward Witten



# From real to complex – Analytic Continuation

## Why can we conduct analytic continuation?

- Tree level scattering amplitudes are rational functions of Lorentz invariants, such as  $p_{i\mu}p_j^\mu$ ,  $p_{i\mu}\epsilon_j^\mu$ .
- **Locality** tells us that any pole of a tree-level amplitude must correspond to a on-shell propagating particle.
- There's only single pole, no branch cuts (logs, square roots, etc) at tree level.



Amplitudes can be shifted to complex plane

# Momentum Shift in BCFW

## What did BCFW do to make the shift?

Here we consider the case in which all particles are massless,  $p_i^2 = 0$  for all  $i = 1, 2, \dots, n$ . Then introduce  $n$  complex-valued vectors  $r_i^\mu$ .

- (i)  $\sum_{i=1}^n r_i^\mu = 0$ ,
- (ii)  $r_i \cdot r_j = 0$  for all  $i, j = 1, 2, \dots, n$ . In particular  $r_i^2 = 0$ ,
- (iii)  $p_i \cdot r_i = 0$  for each  $i$  (no sum).

These vectors  $r_i$  are used to define  $n$  shifted momenta

$$\hat{p}_i^\mu \equiv p_i^\mu + z r_i^\mu \quad \text{with } z \in \mathbb{C}$$



Note that,

- (A) By property (i), momentum conservation holds for the shifted momenta:  $\sum_{i=1}^n \hat{p}_i^\mu = 0$ ,
- (B) By (ii) and (iii), we have  $\hat{p}_i^2 = 0$ , so each shifted momentum is on-shell,
- (C) For a non-trivial subset of generic momenta  $\{p_i\}_{i \in I}$ , define  $P_I^\mu = \sum_{i \in I} p_i^\mu$ .

Then,  $\hat{P}_I^2$  is **linear** in  $z$ :

$$\hat{P}_I^2 = \left( \sum_{i \in I} \hat{p}_i \right)^2 = P_I^2 + 2z P_I \cdot R_I \quad \text{with} \quad R_I = \sum_{i \in I} r_i,$$

because the  $z^2$  term vanishes by property (ii). We can write

$$\hat{P}_I^2 = -\frac{P_I^2}{z_I} (z - z_I) \quad \text{with} \quad z_I = -\frac{P_I^2}{2P_I \cdot R_I}$$

# Fantastic result from Cauchy Theorem

As a result of (A) and (B) (momentum conservation and on-shell), we can consider amplitude  $A_n$  in terms of shifted momentum  $\hat{p}_i^\mu$  instead of original real momentum.

$$A_n \longrightarrow \hat{A}_n(z)$$

and we have known the possible positions of single poles,  $z_I$ , different propagators give us different single poles in the  $z$ -plane.

If we consider the meromorphic function  $\frac{\hat{A}_n(z)}{z}$  in the complex plane, pick a contour that surrounds the simple pole at the origin. ★ The most important point here is that

$$\text{Res}|_{z=0} \frac{\hat{A}_n(z)}{z} = \hat{A}_n(0) = A_n$$

It means that the original amplitude equals to the residue at origin.

From Cauchy Theorem, we can obtain

$$A_n = - \sum_{z_I} \text{Res}|_{z=z_I} \frac{\hat{A}_n(z)}{z} + B_n,$$

where  $B_n$  is the residue of the pole at  $z = \infty$ , called boundary term.

Then, at a  $z_I$  pole, the propagator  $\hat{P}_I^2$  goes to on-shell. In that limit, the shifted amplitude **factorizes** into to on-shell parts (Unitarity)

$$\hat{A}_n(z) \xrightarrow{z \text{ near } z_I} \hat{A}_L(z_I) \frac{1}{\hat{P}_I^2} \hat{A}_R(z_I) = - \frac{z_I}{z - z_I} \hat{A}_L(z_I) \frac{1}{P_I^2} \hat{A}_R(z_I)$$

This makes it easy to evaluate the residue at  $z = z_I$

$$- \text{Res}|_{z=z_I} \frac{\hat{A}_n(z)}{z} = \hat{A}_L(z_I) \frac{1}{P_I^2} \hat{A}_R(z_I) =$$

# Spinor-helicity formalism

The critical point for Spinor-helicity formalism is that a 4-vector  $p_\mu$  in  $(\frac{1}{2}, \frac{1}{2})$  representation can be equivalently represented by a bispinor in  $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$  representation. Specifically

- Massless Case

$$p_\mu \sigma^\mu = p_{\alpha\dot{\alpha}} = \begin{pmatrix} p^0 - p^3 & -p^1 + ip^2 \\ -p^1 - ip^2 & p^0 + p^3 \end{pmatrix}$$

And  $p_{\alpha\dot{\alpha}}$  can be decomposed into product of two Weyl spinors

$$p_{\alpha\dot{\alpha}} = \lambda_\alpha \tilde{\lambda}_{\dot{\alpha}} = |\lambda\rangle[\lambda]$$

Here we can find that there is a ambiguity for the definition, the momentum is invariant under the following redefinition

$$\lambda \rightarrow t^{-1}\lambda, \quad \tilde{\lambda} \rightarrow t\tilde{\lambda}, \quad t \in \mathbb{C}^*$$

same as

$$|\lambda\rangle \rightarrow t^{-1}|\lambda\rangle, \quad [\lambda] \rightarrow t[\lambda]$$

This redefinition just matches the  $U(1)$  little group transformation. The scattering amplitudes should transform covariantly under little group scaling:

$$\mathcal{A}_n(\{|1\rangle, |1], h_1\}, \dots \{t_i^{-1}|i\rangle, t_i|i], h_i\}, \dots) = t_i^{2h_i} \mathcal{A}_n$$

- **Massive Case**

It can also be handled in terms of spinor-helicity variable, see also arXiv:1709.04891 [hep-th] (Nima Arkani-Hamed, Tzu-Chen Huang, Yu-tin Huang).

## 3-point can be completely determined

- 3-particle special kinematics determines an on-shell 3-point amplitude with massless particles depends only on either angle or square brackets of the external momenta.

Let us suppose that it depends only on angle brackets, then we can write down the general ansatz

$$A_3(1^{h_1}, 2^{h_2}, 3^{h_3}) = c \langle 12 \rangle^{x_{12}} \langle 13 \rangle^{x_{13}} \langle 23 \rangle^{x_{23}},$$

Little group scaling tells us that

$$t_1^{2h_1} A_3(1^{h_1}, 2^{h_2}, 3^{h_3}) = c t_1^{-x_{12}} t_1^{-x_{13}} \langle 12 \rangle^{x_{12}} \langle 13 \rangle^{x_{13}} \langle 23 \rangle^{x_{23}}.$$

We can obtain

$$2h_1 = -x_{12} - x_{13}$$

Similarly, we can also obtain

$$2h_2 = -x_{12} - x_{23}, \quad 2h_3 = -x_{13} - x_{23}.$$

Then all index can be solved from this system of equations, so that

$$A_3^{h_1 h_2 h_3} = c \langle 12 \rangle^{h_3 - h_1 - h_2} \langle 31 \rangle^{h_2 - h_1 - h_3} \langle 23 \rangle^{h_1 - h_2 - h_3} \quad h_1 + h_2 + h_3 < 0.$$

Also for square brackets case, we can obtain

$$A_3^{h_1 h_2 h_3} = c' [12]^{h_1 + h_2 - h_3} [23]^{h_2 + h_3 - h_1} [31]^{h_3 + h_1 - h_2} \quad h_1 + h_2 + h_3 > 0.$$

**Example:** 3-gluon amplitude

$$A_3(g_1^-, g_2^-, g_3^+) = g \frac{\langle 12 \rangle^3}{\langle 23 \rangle \langle 31 \rangle}$$

# Contents

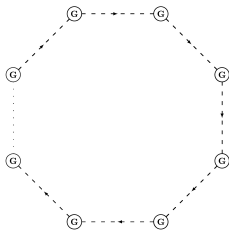
- 1 Why we need new method?
- 2 Preliminary
- 3 Introduction of quiver gauge theory**
- 4 Scattering amplitudes from BCFW



# Introduction of Quiver of Moose gauge theory

- Quiver: A container for carrying arrows
- Moose: A kind of deer with large horns

In the language of field theories, quiver gauge theories contain gauge fields and fermions, summarized in a pictorial representation.



Moose diagram

$N$  – sided polygon

$G$  : gauge group  $SU(m)$

$\rightarrow$ : Unitary scalar fields  $\Phi_{ij}$

# Why we focus on quiver gauge theory?

The lagrangian can be written like

$$\mathcal{L} = - \sum_{i=1}^N \frac{1}{2} \text{Tr}(F_i)^2 + \sum_{i=1}^k \text{Tr}[(D_\mu \Phi_i)^\dagger (D^\mu \Phi_i)],$$

here  $F_i$  refers to the  $i$ th gauge field strength, scalar field  $\Phi_i$  transformed under the **bi-fundamental** representation and the covariant derivative equals to

$$D_\mu \Phi_i = \partial_\mu \Phi_i - i g_i V_{i\mu} \Phi_i + i g_{i+1} \Phi_i V_{i+1\mu}.$$

Here, gauge field and scalar field transformed like

$$\mathbf{A}_{i\mu} \rightarrow g_i(x) \mathbf{A}_{i\mu} g_i^{-1}(x), \quad \Phi_i \rightarrow g_i(x) \Phi_i g_{i+1}^{-1}(x)$$

It is easy to confirm that this theory is invariant under  $\prod_1^N SU(M)$  gauge group.

It has been proposed that this model actually discretized a five-dimension gauge theory with gauge group  $SU(m)$ , where the only the fifth dimension are latticed.

- If  $SU(m)_1$  and  $SU(m)_N$  are connected  $\longrightarrow S^2$  compactification
- If not connected  $\longrightarrow$  Interval compactification

After higgsing, we can obtain a spectrum

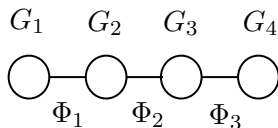
$$M_k^2 = 4g^2 g_s^2 \sin^2 \left( \frac{\pi k}{N} \right)$$

This is precisely the **Kaluza-Klein** spectrum under  $S^2$  compactification.

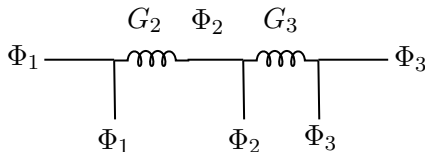
# What is relation to scattering amplitude?

The critical point is **locality**.

- Space-Time Locality  $\rightarrow$  local field theories
- Theory Space Locality  $\rightarrow$  Discretized theory space



If we change this to a scattering diagram



$$\sim 1/z^{\textcircled{4}}$$

# Classification of scattering amplitude

For the simplicity, we started from the two-site gauge theory, means that there are only two gauge fields  $V_1$ ,  $V_2$  and one scalar field  $\Phi$ ,  $\Phi^\dagger$ . The amplitudes in this theory can be classified to many types

- 3-point

- $V_1 \Phi \Phi^\dagger$
- $V_2 \Phi \Phi^\dagger$
- $V_1 V_1 V_1$
- $V_2 V_2 V_2$

- 4-point

- $V_1 V_1 V_1 V_1$
- $V_2 V_2 V_2 V_2$
- $\Phi^\dagger V_1 V_1 \Phi$
- $\Phi V_2 V_2 \Phi^\dagger$
- $\Phi V_2 \Phi^\dagger V_1$
- $\Phi \Phi^\dagger \Phi \Phi^\dagger$

- 5-point

- $V_1 V_1 V_1 V_1 V_1$
- $V_2 V_2 V_2 V_2 V_2$
- $\Phi^\dagger V_1 V_1 V_1 \Phi$
- $\Phi V_2 V_2 V_2 \Phi^\dagger$
- $V_2 \Phi^\dagger V_1 V_1 \Phi$
- $\Phi V_2 V_2 \Phi^\dagger V_1$
- $\Phi \Phi^\dagger \Phi \Phi^\dagger V_1$  (Not yet solved)
- $\Phi \Phi^\dagger \Phi \Phi^\dagger V_2$  (Not yet solved)

- 6-point

- $V_1 V_1 V_1 V_1 V_1 V_1$
- $V_2 V_2 V_2 V_2 V_2 V_2$
- $\Phi^\dagger V_1 V_1 V_1 V_1 \Phi$
- $\Phi V_2 V_2 V_2 V_2 \Phi^\dagger$
- $V_2 V_2 \Phi^\dagger V_1 V_1 \Phi$
- $\Phi V_2 V_2 \Phi^\dagger V_1 V_1$
- $\vdots$
- $\vdots$
- $\vdots$

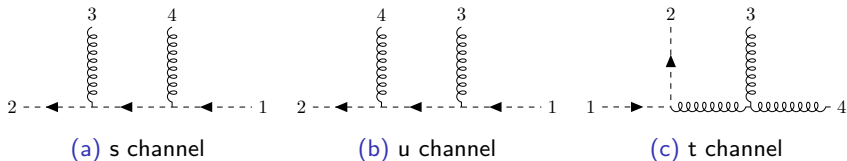
- $nV_1$  or  $nV_2$

This part is completely the same as the pure gluon amplitude, so we can directly borrow the existing results.

$$\text{Parke - Talyor Formula : } A[\cdots, i^-, \cdots, j^-, \cdots] = \frac{\langle ij \rangle^4}{\langle 12 \rangle \langle 23 \rangle \cdots \langle n1 \rangle}$$

Notice that this formula only applies for MHV amplitudes, although the NMHV can be completely solved.

- $\Phi^\dagger V_1 V_1 \Phi$



The color factor can be written respectively as following

$$r_s = \text{Tr}[\Phi_2^\dagger T^{a_3} T^{a_4} \Phi_1], \quad r_u = \text{Tr}[\Phi_2^\dagger T^{a_4} T^{a_3} \Phi_1], \quad r_t = \text{Tr}[\Phi_2^\dagger [T^{a_3}, T^{a_4}] \Phi_1]$$

We can easily obtain a similar Jacobi relation

$$r_t = r_s - r_u$$

Then we can accomplish the color decomposition and define the corresponding color-ordered amplitudes.



For example, in the 4pt. case, the full amplitude can be decomposed to the following form

$$\begin{aligned}\mathcal{A}_4(\Phi^\dagger V_1 V_1 \Phi) &= A_s r_s + A_u r_u + A_t r_t \\ &= A_s r_s + A_u r_u + A_t (r_s - r_u) \\ &= (A_s + A_t) r_s + (A_u - A_t) r_u\end{aligned}$$

The two subamplitudes can be defined as color-ordered amplitude with order  $[1,2,3,4]$  and  $[1,2,4,3]$  respectively.

Of course, for the type  $\Phi^\dagger(nV_1)\Phi$  and  $\Phi(nV_2)\Phi^\dagger$ , we can do the same thing to define the color-ordered amplitudes. It should be noticed that the order only has the relation with the order of external gluon line.

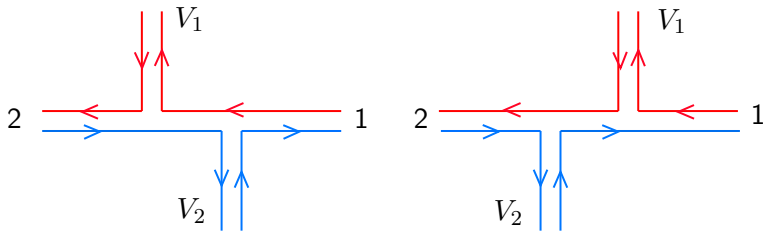
$$[1, 2, \sigma(3), \sigma(4), \cdots, \sigma(n)]$$

- $\Phi V_2 \Phi^\dagger V_1$

The color structure for this kind of amplitude has special form, like

$$(T_1^a)_{ij}(T_2^b)_{\bar{j}\bar{i}}$$

It is more straightforward to observe the color structure in terms of double line notation as follows



# OPP(Order preserving permutation)

From the 4 point case, we have known that the relative order between gauge boson 1 and gauge boson 2 does not affect the color structure. Thus, it is necessary to introduce the **OPP(Order preserving permutation)**. For example:

$$(3_1, 4_1, 5_2) \quad (3_1, 5_2, 4_1) \quad (5_2, 3_1, 4_1)$$

These three permutations are different OPP for  $(3_1, 4_1, 5_2)$ , so that give us the same color factor.

# Contents

- 1 Why we need new method?
- 2 Preliminary
- 3 Introduction of quiver gauge theory
- 4 Scattering amplitudes from BCFW

## Basic building block – 3-point

From the previous section, we have known that there are only two kinds of 3 point amplitude

$$\begin{aligned} A[1, 2, 3^+] &= \frac{[23][31]}{[12]}, & A[1, 2, 3^-] &= \frac{\langle 23 \rangle \langle 31 \rangle}{\langle 12 \rangle} \\ A[3^+, 4^+, 5^-] &= \frac{[34]^3}{[45][53]}, & A[3^-, 4^-, 5_+] &= \frac{\langle 34 \rangle^3}{\langle 45 \rangle \langle 53 \rangle} \end{aligned}$$

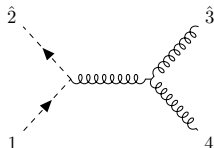
By using the 3 point building block, we can construct 4 point color-ordered amplitudes from BCFW recursion relation.

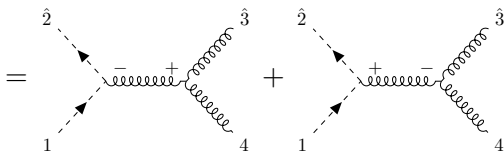
## 4 point from BCFW

•  $\Phi^\dagger V_1 V_1 \Phi$

Here we compute the color-ordered amplitude  $A[1, 2, 3^+, 4^-]$ . We choose  $[2, 3\rangle$  shift

$$\begin{aligned} |\hat{2}\rangle &= |2\rangle - z|3\rangle, & |\hat{2}\rangle &= |2\rangle \\ |\hat{3}\rangle &= |3\rangle, & |\hat{3}\rangle &= |3\rangle + z|2\rangle \end{aligned}$$

$$A[1, 2, 3^+, 4^-] = \sum_h$$




We denote these two different BCFW channels as  $A_1$  and  $A_2$ , then

$$\begin{aligned}
 A_1 &= \frac{\langle \hat{2}\hat{I} \rangle \langle \hat{I}1 \rangle}{\langle 1\hat{I} \rangle} \times \frac{1}{s_{12}} \times \frac{[\hat{I}\hat{3}]^3}{[\hat{3}4][4\hat{I}]} \\
 &= (-1) \frac{\langle 14 \rangle^2 \langle 24 \rangle^2}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle}
 \end{aligned}$$

where we use the fact  $|\hat{2}\rangle = |2\rangle$ ,  $|\hat{3}\rangle = |3\rangle$ , and the **Fierz Identity**

$$[ij][kl] + [il][jk] + [ik][lj] = 0$$

Similarly, we can obtain

$$A[1, 2, 3^-, 4^+] = (-1) \frac{\langle 13 \rangle^2 \langle 23 \rangle^2}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle}$$

★ Bonus relation

$$A[1, 2, 3^+, 4^+] = A[1, 2, 3^-, 4^-] = 0$$