THE INFRARED STRUCTURE OF GAUGE AMPLITUDES IN THE HIGH-ENERGY LIMIT

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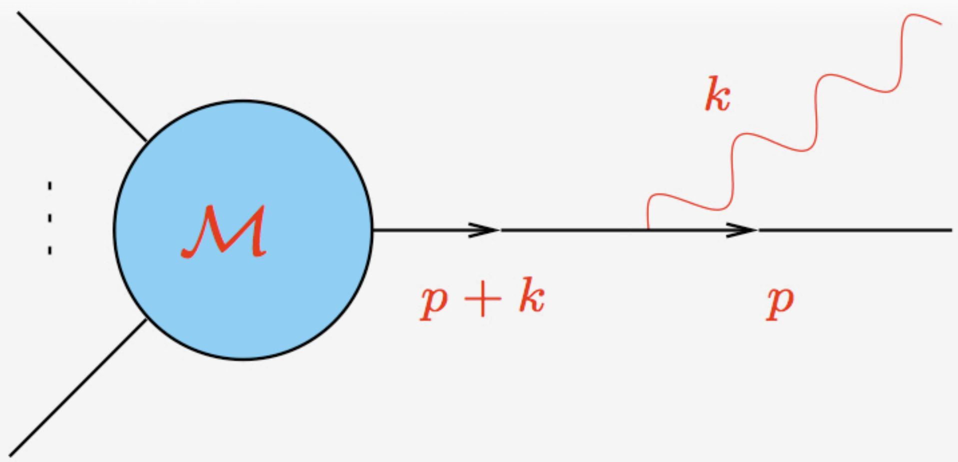
Outline

- Infrared divergences to all orders
- The dipole formula (with E. Gardi)
- The high-energy limit (with V. Del Duca, C. Duhr, E. Gardi, C. White)
- Reggeization and beyond
- Outlook

ON INFRARED DIVERGENCES

Textbook theory ...



$$\to \quad -ig\bar{u}(p)\not{\epsilon}(k)t_a \frac{i(\not{p} + \not{k})}{(p+k)^2 + i\varepsilon}\mathcal{M}\,,$$

Singularities arise only when propagators go on shell

$$2p \cdot k = 2p_0 k_0 (1 - \cos\theta_{pk}) = 0\,,$$
$$\to k_0 = 0 \ (IR); \quad \cos\theta_{pk} = 1\,.$$

➡ Emission is not suppressed at long distances

➡ Isolated charged particles are not true asymptotic states of unbroken gauge theories

- A serious problem: the S matrix does not exist in the usual Fock space
- Possible solutions: construct finite transition probabilities (KLN theorem)
  construct better asymptotic states (coherent states)
- Long-distance singularities obey a pattern of exponentiation

$$\mathcal{M} = \mathcal{M}_0 \left[1 - \kappa \frac{\alpha}{\pi} \frac{1}{\epsilon} + \ldots \right] \quad \Rightarrow \quad \mathcal{M} = \mathcal{M}_0 \exp\left[-\kappa \frac{\alpha}{\pi} \frac{1}{\epsilon} + \ldots \right]$$

… and Practice

Just a formal issue in Quantum Field Theory? Are there practical applications?

- Higher order QCD calculations at colliders hinge upon cancellation of divergences between virtual corrections and real emission contributions
  - Cancellation must be performed analytically before numerical integrations
  - Need local counterterms for matrix elements in all singular regions
  - State of the art: NLO multileg, NNLO for (some) color-singlet processes

- Cancellations leave behind large logarithms: they must be resummed

$$\underbrace{\frac{1}{\epsilon}}_{\text{virtual}} + \underbrace{(Q^2)^\epsilon \int_0^{m^2} \frac{dk^2}{(k^2)^{1+\epsilon}}}_{\text{real}} \qquad \Longrightarrow \qquad \ln(m^2/Q^2)$$

  - For inclusive observables: analytic resummation to high logarithmic accuracy.
  - For exclusive final states: parton shower event generators, (N)LL accuracy.

- Resummation probes the all-order structure of perturbation theory
  - Power-suppressed corrections to QCD cross sections can be studied
  - Links to the strong coupling regime can be established for SUSY gauge theories.

TOOLS

Dimensional regularization

Exponentiation of infrared poles requires solving d-dimensional evolution equations. The running coupling in d = 4 - 2 ε obeys

$$\mu \frac{\partial \overline{\alpha}}{\partial \mu} \equiv \beta(\epsilon, \overline{\alpha}) = -2\,\epsilon\,\overline{\alpha} + \hat{\beta}(\overline{\alpha}) \quad , \quad \hat{\beta}(\overline{\alpha}) = -\frac{\overline{\alpha}^2}{2\pi} \sum_{n=0}^{\infty} b_n \left(\frac{\overline{\alpha}}{\pi} \right)^n \quad .$$

The one-loop solution is

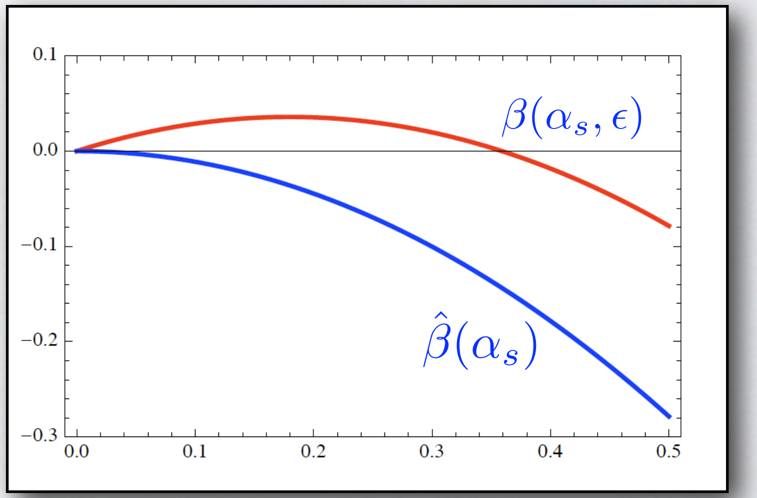
$$\overline{\alpha}\left(\mu^2, \epsilon\right) = \alpha_s(\mu_0^2) \left[\left(\frac{\mu^2}{\mu_0^2} \right)^{\epsilon} - \frac{1}{\epsilon} \left(1 - \left(\frac{\mu^2}{\mu_0^2} \right)^{\epsilon} \right) \frac{b_0}{4\pi} \alpha_s(\mu_0^2) \right]^{-1} \quad .$$

The β function develops an IR-free fixed point, so that the coupling vanishes at μ = 0 for fixed ε < 0. The Landau pole is at

$$\mu^2 = \Lambda^2 \equiv Q^2 \left(1 + \frac{4\pi\epsilon}{b_0 \alpha_s(Q^2)} \right)^{-1/\epsilon} \quad .$$

➡ Integrations over the scale of the coupling can be analytically performed.

➡ All infrared and collinear poles arise by integration over the scale of the running coupling.



For negative **ε** the beta function develops a second zero, O(**ε**) from the origin.

Factorization

All factorizations separating dynamics at different energy scales lead to resummation of logarithms of the ratio of scales.

A textbook example is collinear factorization for DIS structure functions.

- Collinear factorization separates the dependence on the physical scale Q^2 from the dependence on collinear cutoffs (parton masses m^2). For Mellin moments one gets

$$\widetilde{F}_2\left(N, \frac{Q^2}{m^2}, \alpha_s\right) = \widetilde{C}\left(N, \frac{Q^2}{\mu_F^2}, \alpha_s\right) \widetilde{f}\left(N, \frac{\mu_F^2}{m^2}, \alpha_s\right) .$$

- Factorization requires the introduction of an arbitrarily chosen scale μ_F. Results must be independent of the arbitrary choice of μ_F.

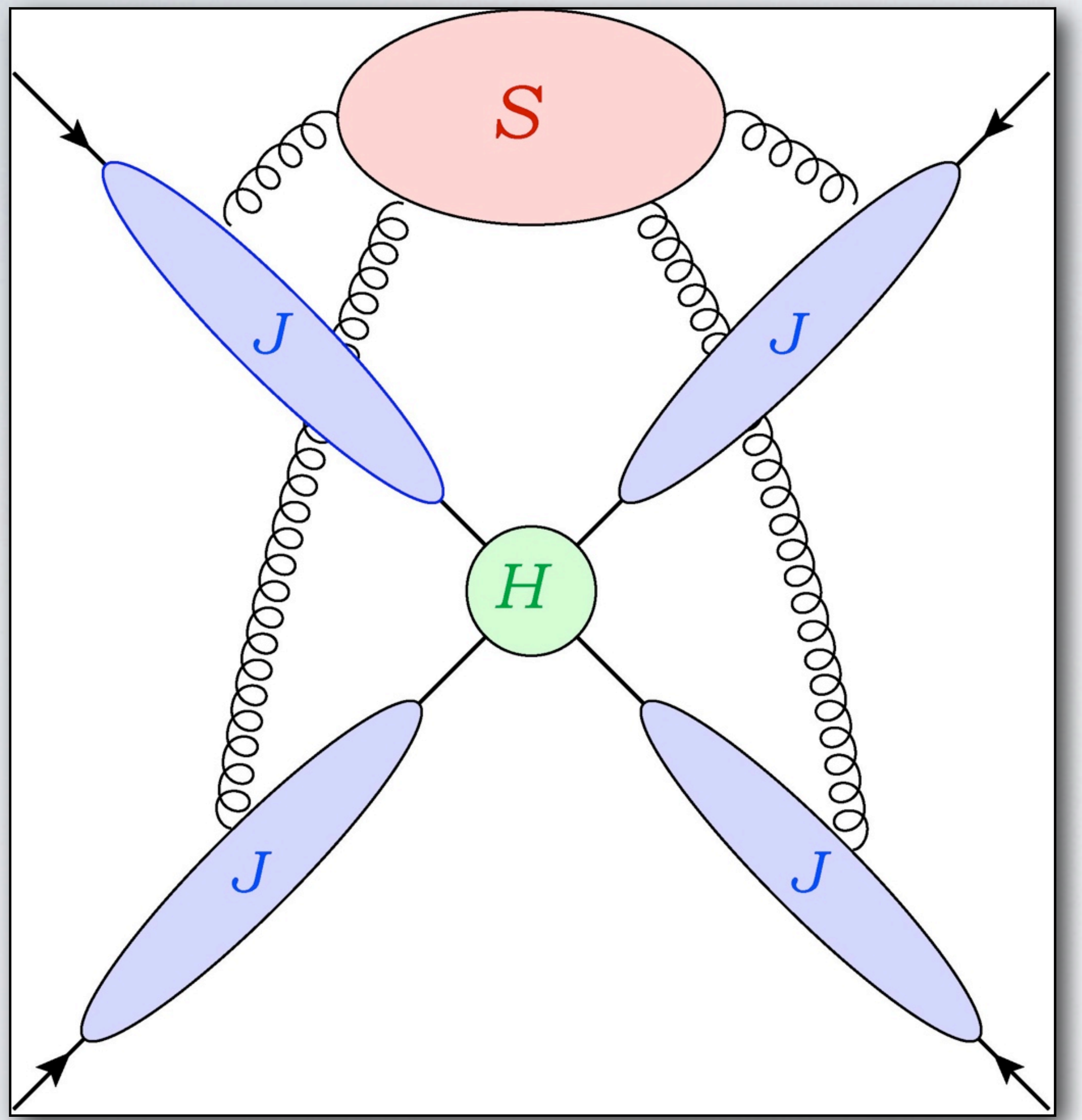
$$\frac{d\widetilde{F}_2}{d\mu_F} = 0 \quad \longrightarrow \quad \frac{d \log \widetilde{f}}{d \log \mu_F} = \gamma_N\left(\alpha_s\right) .$$

- The simple functional dependence of the factors is dictated by separation of variables.

- Proving factorization is the difficult step: it requires all-order diagrammatic analyses, or OPE. Evolution equations for parton distributions follow automatically.

- Solving Altarelli-Parisi evolution resums logarithms of Q^2/μ_F^2 into evolved parton distributions (or fragmentation functions).

Sudakov Factorization



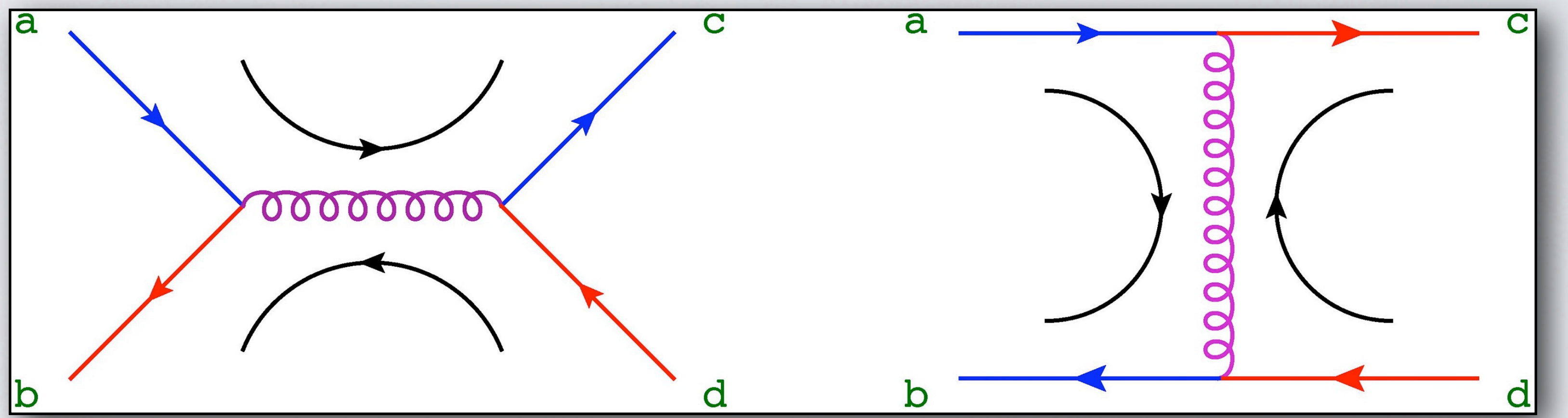
Leading integration regions in loop momentum space for Sudakov factorization

- Sudakov logarithms are remainders of infrared and collinear divergences.
- Divergences arise in scattering amplitudes from leading regions in loop momentum space.
- Soft gluons factorize both form hard (easy) and from collinear (intricate) virtual exchanges.
- Jet functions J represent color singlet evolution of external hard partons.
- The soft function S is a matrix mixing the available color representations.
- In the planar limit soft exchanges are confined to wedges: S is proportional to the identity.
- In the planar limit S can be reabsorbed defining jets as square roots of elementary form factors.
- Beyond the planar limit S is determined by an anomalous dimension matrix Γ_S.

Color flow

In order to understand the matrix structure of the soft function it is sufficient to consider the simple case of quark-antiquark scattering.

At tree level



Tree-level diagrams and color flows for quark-antiquark scattering

For this process only two color structures are possible. A basis in the space of available color tensors is

$$c_{abcd}^{(1)} = \delta_{ab}\delta_{cd}\,, \qquad c_{abcd}^{(2)} = \delta_{ac}\delta_{bd}$$

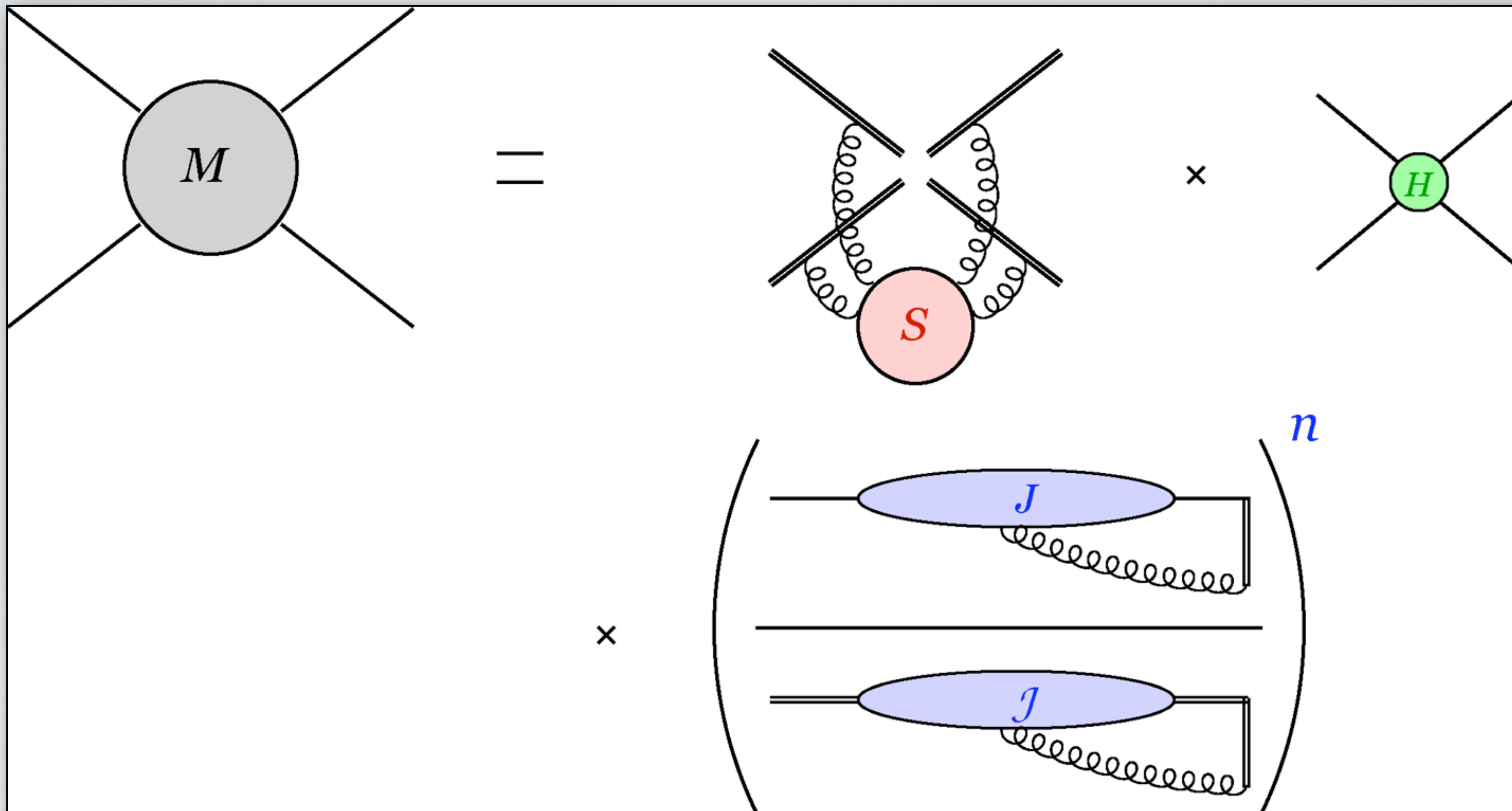
The matrix element is a vector in this space, and the Born cross section is

$$\mathcal{M}_{abcd} = \mathcal{M}_1\, c_{abcd}^{(1)} + \mathcal{M}_2\, c_{abcd}^{(2)} \longrightarrow \sum_{\text{color}} |\mathcal{M}|^2 = \sum_{J,L} \mathcal{M}_J\, \mathcal{M}_L^*\ \text{tr}\left[c_{abcd}^{(J)} \left(c_{abcd}^{(L)}\right)^\dagger\right] \equiv \text{Tr}\Big[HS\Big]_0$$

A virtual soft gluon will reshuffle color and mix the components of this vector

$$\text{QED}: \quad \mathcal{M}_{\text{div}} = S_{\text{div}} \mathcal{M}_{\text{Born}}\,; \qquad \text{QCD}: \quad [\mathcal{M}_{\text{div}}]_J = [S_{\text{div}}]_{JL}\, [\mathcal{M}_{\text{Born}}]_L$$

Sudakov factorization: pictorial



A pictorial representation of Sudakov factorization for fixed-angle scattering amplitudes

Operator Definitions

The precise functional form of this graphical factorization is

$$\mathcal{M}_L\left(p_i/\mu, \alpha_s(\mu^2), \epsilon\right) = \mathcal{S}_{LK}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right) H_K\left(\frac{p_i \cdot p_j}{\mu^2}, \frac{(p_i \cdot n_i)^2}{n_i^2 \mu^2}, \alpha_s(\mu^2)\right)$$
$$\times \prod_{i=1}^{n} \left[J_i\left(\frac{(p_i \cdot n_i)^2}{n_i^2 \mu^2}, \alpha_s(\mu^2), \epsilon\right) \Big/ \mathcal{J}_i\left(\frac{(\beta_i \cdot n_i)^2}{n_i^2}, \alpha_s(\mu^2), \epsilon\right)\right] ,$$

We introduced factorization vectors n_i^μ , $n_i^2 \neq 0$ to define the jets,

$$J\left(\frac{(p \cdot n)^2}{n^2 \mu^2}, \alpha_s(\mu^2), \epsilon\right) u(p) = \langle 0 \,|\Phi_n(\infty, 0)\, \psi(0)\,|p\rangle \,.$$

where Φ_n is the Wilson line operator along the direction n^μ ,

$$\Phi_n(\lambda_2, \lambda_1) = P \exp\left[ig \int_{\lambda_1}^{\lambda_2} d\lambda\, n \cdot A(\lambda n)\right] \,.$$

The vectors n^μ :

- Ensure gauge invariance of the jets.
- Separate collinear gluons from wide-angle soft ones.
- Replace other hard partons with a collinear-safe absorber.

Soft Matrices

The soft function S is a matrix, mixing the available color tensors. It is defined by a correlator of Wilson lines.

$$(c_L)_{\{\alpha_k\}} \mathcal{S}_{LK}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right) = \sum_{\{\eta_k\}} \langle 0| \prod_{i=1}^{n} \left[\Phi_{\beta_i}(\infty, 0)_{\alpha_k, \eta_k}\right] |0\rangle \, (c_K)_{\{\eta_k\}} \, ,$$

The soft function S obeys a matrix RG evolution equation

$$\mu \frac{d}{d\mu} \mathcal{S}_{IK}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right) = - \mathcal{S}_{IJ}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right) \Gamma^{\mathcal{S}}_{JK}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right)$$

- Γ^S is singular due to overlapping UV and collinear poles.

S is a pure counterterm. In dimensional regularization, using $\alpha_s(\mu^2 = 0, \varepsilon < 0) = 0$,

$$\mathcal{S}\left(\beta_i \cdot \beta_j, \alpha_s(\mu^2), \epsilon\right) = P \exp\left[-\frac{1}{2} \int_0^{\mu^2} \frac{d\xi^2}{\xi^2} \Gamma^{\mathcal{S}}\left(\beta_i \cdot \beta_j, \alpha_s(\xi^2, \epsilon), \epsilon\right)\right] .$$

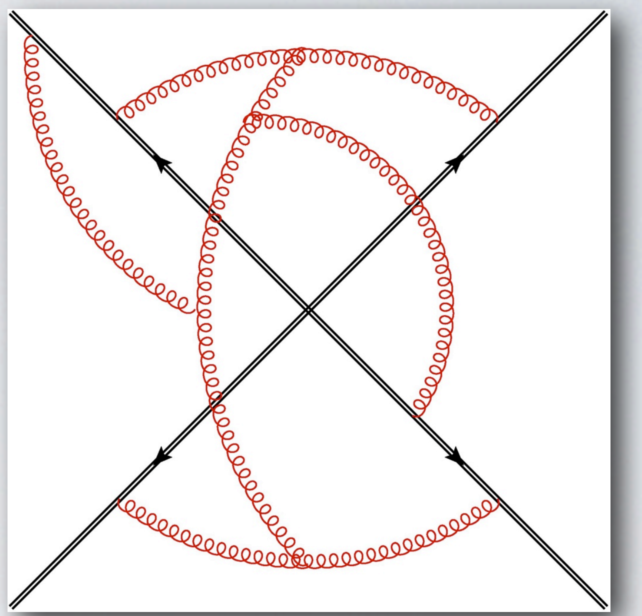
The determination of the soft anomalous dimension matrix Γ^S is the keystone of the resummation program for multiparton amplitudes and cross sections.

- It governs the interplay of color exchange with kinematics in multiparton processes.
- It is the only source of multiparton correlations for singular contributions.
- Collinear effects are \`color singlet' and can be extracted from two-parton scatterings.

THE DIPOLE FORMULA

Surprising Simplicity

- The matrix Γ_S can be computed from the UV poles of S.
- Computations can be performed directly for the exponent: the relevant diagrams are called "webs".
- Γ_S appears highly complex at high orders.
- g-loop webs directly correlate color and kinematics of up to g+1 Wilson lines.



A web contributing to the soft anomalous dimension matrix

The two-loop calculation (Aybat, Dixon, Sterman 06) leads to a surprising result: for any number of external massless partons

$$\Gamma_{\mathcal{S}}^{(2)} = \frac{\kappa}{2}\,\Gamma_{\mathcal{S}}^{(1)} \qquad \kappa = \left(\frac{67}{18} - \zeta(2)\right) C_A - \frac{10}{9} T_F C_F \,.$$

➡ No new kinematic dependence; no new matrix structure.

➡ κ is the two-loop coefficient of $\gamma_K(\alpha_s)$, rescaled by the appropriate quadratic Casimir,

$$\gamma_K^{(i)}(\alpha_s) = C^{(i)}\left[2\,\frac{\alpha_s}{\pi} + \kappa\left(\frac{\alpha_s}{\pi}\right)^2 + \mathcal{O}\left(\alpha_s^3\right)\right] \,.$$

The Dipole Formula

The two-loop result led to an all-order understanding. For massless partons, the soft matrix obeys a set of exact equations that correlate color exchange with kinematics.

The simplest solution to these equations is a sum over color dipoles (Becher, Neubert; Gardi, LM, 09). It leads to an ansatz for the all-order singularity structure of all multiparton fixed-angle massless scattering amplitudes: the dipole formula.

- All soft and collinear singularities can be collected in a multiplicative operator Z

$$\mathcal{M}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right) = Z\left(\frac{p_i}{\mu_f}, \alpha_s(\mu_f^2), \epsilon\right)\ \mathcal{H}\left(\frac{p_i}{\mu}, \frac{\mu_f}{\mu}, \alpha_s(\mu^2), \epsilon\right) ,$$

- Z contains both soft singularities from S, and collinear ones from the jet functions. It must satisfy its own matrix RG equation

$$\frac{d}{d\ln\mu} Z\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right) = -Z\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right) \Gamma\left(\frac{p_i}{\mu}, \alpha_s(\mu^2)\right) .$$

The matrix Γ inherits the dipole structure from the soft matrix. It reads

$$\boxed{\Gamma_{\rm dip}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2)\right) = -\frac{1}{4}\widehat{\gamma}_K\left(\alpha_s(\mu^2)\right)\sum_{j\neq i}\ln\left(\frac{-2\, p_i\cdot p_j}{\mu^2}\right)\mathbf{T}_i\cdot\mathbf{T}_j + \sum_{i=1}^{n}\gamma_{J_i}\left(\alpha_s(\mu^2)\right) .}$$

Note that all singularities are generated by integration over the scale of the coupling.

Features of the dipole formula

- All known results for IR divergences of massless gauge theory amplitudes are recovered.
- The absence of multiparton correlations implies remarkable diagrammatic cancellations.
- The color matrix structure is fixed at one loop: path-ordering is not needed.
- All divergences are determined by a handful of anomalous dimensions.
- The cusp anomalous dimension plays a very special role: a universal IR coupling.

Can this be the definitive answer for IR divergences in massless non-abelian gauge theories?

- There are precisely two sources of possible corrections.
  - Quadrupole correlations may enter starting at three loops: they must be tightly constrained functions of conformal cross ratios of parton momenta.

$$\Gamma\left(\frac{p_i}{\mu}, \alpha_s(\mu^2)\right) = \Gamma_{\rm dip}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2)\right) + \Delta\left(\rho_{ijkl}, \alpha_s(\mu^2)\right) \, , \qquad \rho_{ijkl} = \frac{p_i \cdot p_j \, p_k \cdot p_l}{p_i \cdot p_k \, p_j \cdot p_l}$$

  - The cusp anomalous dimension may violate Casimir scaling beyond three loops.

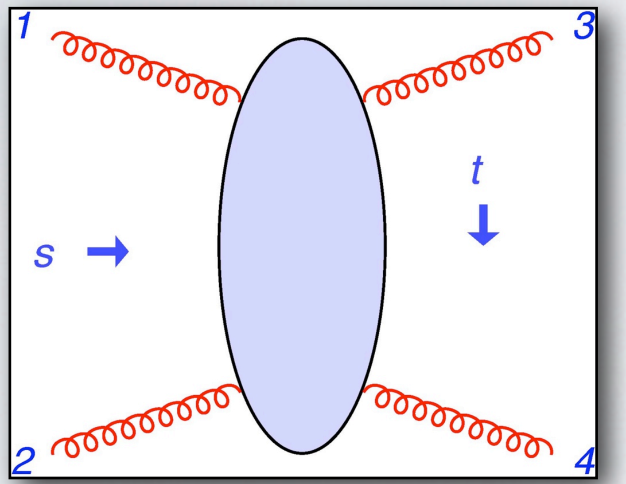
$$\gamma_K^{(i)}(\alpha_s) = C_i \, \widehat{\gamma}_K(\alpha_s) + \widetilde{\gamma}_K^{(i)}(\alpha_s)$$

  - The functional form of Δ is further constrained by: collinear limits, Bose symmetry and transcendentality bounds (Becher, Neubert; Dixon, Gardi, LM, 09).
  - A four-loop analysis indicates that Casimir scaling holds (Becher, Neubert, Vernazza).

THE HIGH-ENERGY LIMIT

Reggeization

- Studies of the high-energy limit predate the modern era of quantum field theory.
- In the t/s → 0 limit particles exchanged in the t-channel (may) `Reggeize'.



Gluon-gluon scattering: the t-channel gluon Reggeizes

- Large logarithms of s/t are generated by a simple replacement of the t-channel propagator,

$$\frac{1}{t} \longrightarrow \frac{1}{t}\left(\frac{s}{-t}\right)^{\alpha(t)}$$

- The Regge trajectory has a perturbative expansion, with IR divergent coefficients

$$\alpha(t) = \frac{\alpha_s(-t,\epsilon)}{4\pi}\,\alpha^{(1)} + \left(\frac{\alpha_s(-t,\epsilon)}{4\pi}\right)^2 \alpha^{(2)} + \mathcal{O}\left(\alpha_s^3\right)$$

- The gluon has been shown to Reggeize at NLL, and the two-loop Regge trajectory is known.
- For example, for gluon-gluon scattering the matrix element obeys Regge factorization

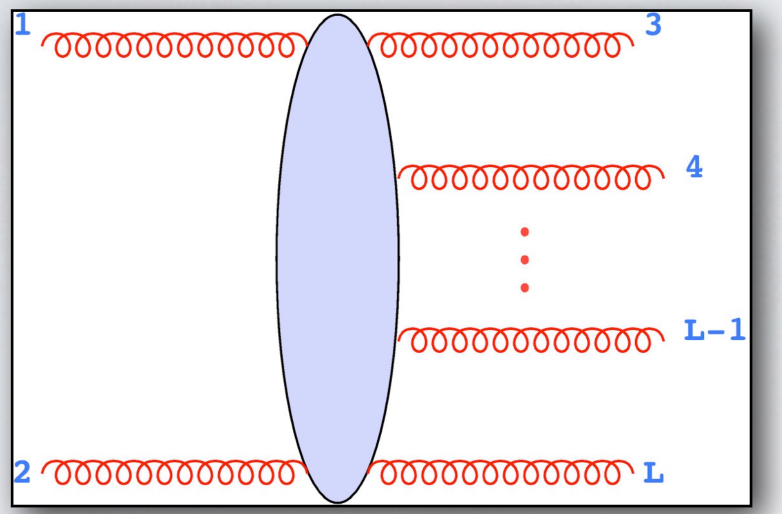
$$\mathcal{M}^{gg\to gg}_{a_1a_2a_3a_4}(s,t) = 2\,g_s^2\,\frac{s}{t}\left[(T^b)_{a_1a_3}C_{\lambda_1\lambda_3}(k_1,k_3)\right]\left(\frac{s}{-t}\right)^{\alpha(t)}\left[(T_b)_{a_2a_4}C_{\lambda_2\lambda_4}(k_2,k_4)\right]$$

with the perturbative coefficients

$$\alpha^{(1)} = C_A\,\frac{\widehat{\gamma}_K^{(1)}}{\epsilon} = C_A\,\frac{2}{\epsilon} \qquad \alpha^{(2)} = C_A\left[-\frac{b_0}{\epsilon^2} + \widehat{\gamma}_K^{(2)}\,\frac{2}{\epsilon} + C_A\left(\frac{404}{27} - 2\zeta_3\right) + n_f\left(-\frac{56}{27}\right)\right]$$

Multi-Regge kinematics

- Reggeization follows form the dominance of t-channel ladder diagrams as $t/s \to 0$.
- By unitarity, multi-gluon emission must similarly simplify in the high-energy limit.
- Regge factorization extends to multi-particle emission in `Multi-Regge' kinematics.



Multi-gluon emission and Multi-Regge kinematics

$$y_3 \gg y_4 \gg \ldots \gg y_L \,, \qquad |k_i^\perp| \simeq |k_j^\perp| \,, \qquad \forall i, j$$

$$-s \equiv -s_{12} \simeq |k_3^\perp| \, |k_L^\perp| \, \mathrm{e}^{y_3 - y_L} \, \mathrm{e}^{\mathrm{i}\pi}$$

$$-s_{ij} \simeq |k_i^\perp| \, |k_j^\perp| \, \mathrm{e}^{y_i - y_j} \, \mathrm{e}^{\mathrm{i}\pi} \,, \quad 3 \leq i < j \leq L$$

- Large logarithms of s/t_i are generated by the Reggeization of t-channel propagators, as

$$\begin{aligned} \mathcal{M}^{gg \to (L-2)g}_{a_1 \ldots a_L} \;=\; & 2 \, g_s^3 \, s \left[(T^b)_{a_1 a_3} \, C_{\lambda_1 \lambda_3}(k_1, k_3) \right] \left[\frac{1}{t_1} \left(\frac{s_{34}}{-t_1} \right)^{\alpha(t_1)} \right] \\ & \times \left[(T^{a_4})_{bc} \, V_{\lambda_4}(q_1, q_2) \right] \left[\frac{1}{t_2} \left(\frac{s_{45}}{-t_2} \right)^{\alpha(t_2)} \right] \ldots \left[(T^c)_{a_2 a_L} \, C_{\lambda_2 \lambda_L}(k_2, k_L) \right] \end{aligned}$$

- The impact factors C and the Lipatov vertices V are universal and independent of s.

The dipole formula at high energy

- Introducing Mandelstam color operators, and using color and momentum conservation

$$\begin{aligned} \mathbf{T}_s &= \mathbf{T}_1 + \mathbf{T}_2 = -(\mathbf{T}_3 + \mathbf{T}_4) \, , \\ \mathbf{T}_t &= \mathbf{T}_1 + \mathbf{T}_3 = -(\mathbf{T}_2 + \mathbf{T}_4) \, , \\ \mathbf{T}_u &= \mathbf{T}_1 + \mathbf{T}_4 = -(\mathbf{T}_2 + \mathbf{T}_3) \end{aligned}$$

$$s + t + u = 0$$

$$\mathbf{T}_s^2 + \mathbf{T}_t^2 + \mathbf{T}_u^2 = \sum_{i=1}^{4} C_i$$

it is easy to see that the infrared dipole operator Z factorizes in the high-energy limit

$$Z\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right) = \widetilde{Z}\left(\frac{s}{t}, \alpha_s(\mu^2), \epsilon\right) \, Z_{\mathbf{1}}\left(\frac{t}{\mu^2}, \alpha_s(\mu^2), \epsilon\right)$$

- The operator Z_1 is s-independent and proportional to the unit matrix in color space.
- Color dependence and s dependence are collected in the factor

$$\widetilde{Z}\left(\frac{s}{t}, \alpha_s(\mu^2), \epsilon\right) = \exp\left\{ K\Big(\alpha_s(\mu^2), \epsilon\Big) \left[\ln\left(\frac{s}{-t}\right) \mathbf{T}_t^2 + i\pi \, \mathbf{T}_s^2 \right] \right\} ,$$

where the coupling dependence is (once again!) completely determined by the cusp anomalous dimension and by the β function, through

$$K\Big(\alpha_s(\mu^2), \epsilon\Big) \equiv -\frac{1}{4} \int_0^{\mu^2} \frac{d\lambda^2}{\lambda^2} \, \widehat{\gamma}_K \left(\alpha_s(\lambda^2, \epsilon)\right)$$

- The simple structure of the high-energy operator governs Reggeization and its breaking.

Reggeization of leading logarithms

- At leading logarithmic accuracy, the (imaginary) s-channel contribution can be dropped, and the dipole operator becomes diagonal in a t-channel basis.

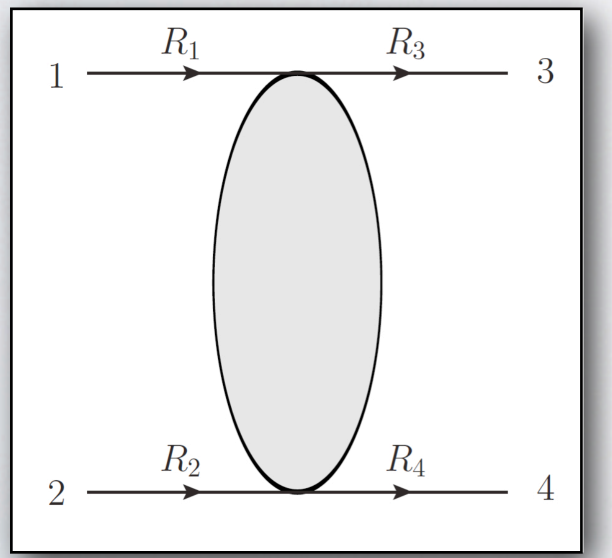
$$\mathcal{M}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right) = \exp\left\{K\Big(\alpha_s(\mu^2), \epsilon\Big) \ln\left(\frac{s}{-t}\right) \mathbf{T}_t^2\right\} Z_1 \, \mathcal{H}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2), \epsilon\right)$$

- If, at LO and at leading power in t/s, the scattering is dominated by t-channel exchange, then the hard function is an eigenstate of the color operator $\mathbf{T}_t^2$

$$\mathbf{T}_t^2 \, \mathcal{H}^{gg\to gg} \xrightarrow{|t/s|\to 0} C_t \, \mathcal{H}_t^{gg\to gg}$$

- Leading-logarithmic Reggeization for arbitrary t-channel color representations follows

$$\mathcal{M}^{gg\to gg} = \left(\frac{s}{-t}\right)^{C_A \, K\left(\alpha_s(\mu^2), \epsilon\right)} Z_1 \, \mathcal{H}_t^{gg\to gg}$$



Scattering for generic color exchange

- The LL Regge trajectory is universal and obeys Casimir scaling.
- Scattering of arbitrary color representations can be analyzed
  Example: let 1 and 2 be antiquarks, 4 a gluon and 3 a sextet; use

$$\overline{\mathbf{3}} \otimes \mathbf{6} = \mathbf{3} \oplus \mathbf{15} \qquad\qquad \overline{\mathbf{3}} \otimes \mathbf{8}_a = \overline{\mathbf{3}} \oplus \mathbf{6} \oplus \overline{\mathbf{15}}$$

LL Reggeization of the **3** and **15** t-channel exchanges follows.

Beyond leading logarithms

- The high-energy infrared operator can be systematically expanded beyond LL, using the Baker-Campbell-Hausdorff formula. At NLL one finds a series of commutators

$$\widetilde{Z}\left(\frac{s}{t},\alpha_s,\epsilon\right)\Big|_{\rm NLL} = \left(\frac{s}{-t}\right)^{K(\alpha_s,\epsilon)\,\mathbf{T}_t^2}\Bigg\{1+{\rm i}\,\pi K\left(\alpha_s,\epsilon\right)\left[\mathbf{T}_s^2-\frac{K\left(\alpha_s,\epsilon\right)}{2!}\ln\left(\frac{s}{-t}\right)\left[\mathbf{T}_t^2,\mathbf{T}_s^2\right]\right.$$
$$\left.+\frac{K^2\left(\alpha_s,\epsilon\right)}{3!}\ln^2\left(\frac{s}{-t}\right)\left[\mathbf{T}_t^2,\left[\mathbf{T}_t^2,\mathbf{T}_s^2\right]\right]+\ldots\right]\Bigg\}$$

- The real part of the amplitude Reggeizes also at NLL for arbitrary t-channel exchanges.
- At NNLL Reggeization generically breaks down also for the real part of the amplitude.
  - At two loops, terms that are non-logarithmic and non-diagonal in a t-channel basis arise

$$\mathcal{E}_0\left(\alpha_s,\epsilon\right)\equiv-\frac{1}{2}\pi^2\,K^2\left(\alpha_s,\epsilon\right)\,\left(\mathbf{T}_s^2\right)^2$$

  - At three loops, the first Reggeization-breaking logarithms of s/t arise, generated by

$$\mathcal{E}_1\left(\frac{s}{t},\alpha_s,\epsilon\right)\equiv-\frac{\pi^2}{3}\,K^3(\alpha_s,\epsilon)\,\ln\left(\frac{s}{-t}\right)\left[\mathbf{T}_s^2,\left[\mathbf{T}_t^2,\mathbf{T}_s^2\right]\right]$$

- **NOTE**
  - In the planar limit ($N_C\to\infty$) all commutators vanish and Reggeization holds also beyond NLL (as perhaps expected from string theory).
  - Possible quadrupole corrections to the dipole formula cannot come to the rescue.

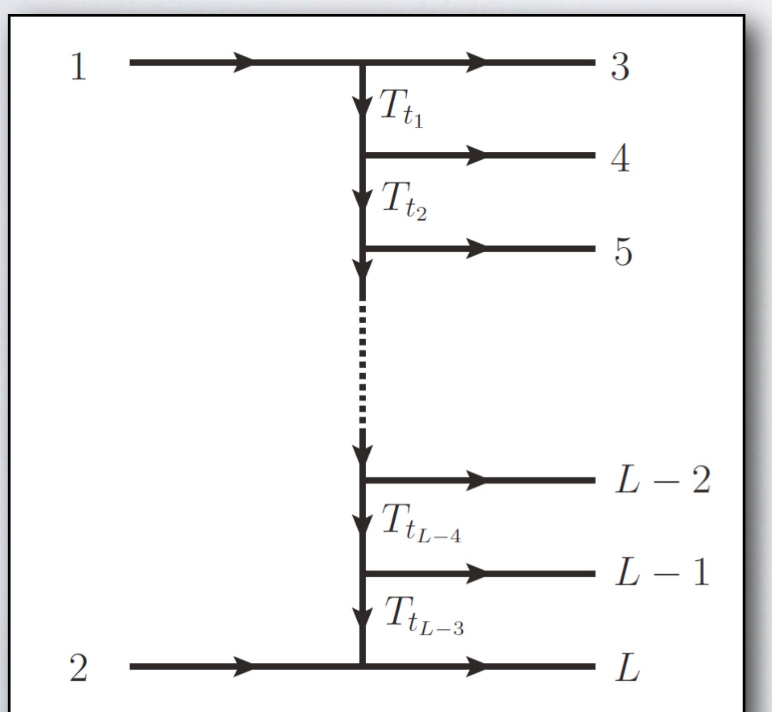
Multi-Regge kinematics

- The dipole formula applies for any number of particles: we expect similar simplifications in Multi-Regge kinematics, and similar results concerning Reggeization.
- Indeed, one can prove recursively that the dipole operator Z factorizes in MR kinematics, as

$$Z\left(\frac{p_l}{\mu}, \alpha_s(\mu^2), \epsilon\right) = \widetilde{Z}^{\rm MR}\Big(\Delta y_k, \alpha_s(\mu^2), \epsilon\Big)\ Z_{\mathbf{1}}^{\rm MR}\left(\frac{|k_i^\perp|}{\mu}, \alpha_s(\mu^2), \epsilon\right)$$

- The Multi-Regge high-energy operator has again a simple structure.

$$\widetilde{Z}^{\rm MR}\Big(\Delta y_k, \alpha_s(\mu^2), \epsilon\Big) = \exp\left\{K\Big(\alpha_s(\mu^2), \epsilon\Big)\left[\sum_{k=3}^{L-1}\mathbf{T}^2_{t_{k-2}}\,\Delta y_k + {\rm i}\pi\mathbf{T}_s^2\right]\right\}$$



Color structure in Multi-Regge kinematics

- We have defined the t-channel color operators

$$\mathbf{T}_{t_k} = \mathbf{T}_1 + \sum_{p=1}^{k}\mathbf{T}_{p+2}$$

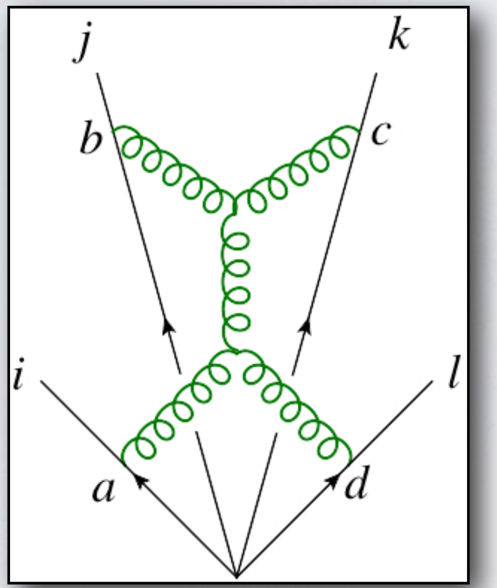
- A t-channel basis of common eigenstates of $\mathbf{T}_{t_k}$ can be constructed using Clebsch-Gordan coefficients.
- The operators $\mathbf{T}_{t_k}$ thus commute, and each color representation contributing to the hard function in the high-energy limit Reggeizes separately at LL.

Constraining quadrupoles

- Known results on the high-energy limit of QCD amplitudes imply new constraints on quadrupole corrections to the dipole formula at three loops and beyond.
- Previous analyses using collinear constraint, Bose symmetry and transcendentality bounds could not exclude a class of correction, including for example

$$\Delta^{(212)}(\rho_{ijkl}, \alpha_s) = \left(\frac{\alpha_s}{\pi}\right)^3 \mathbf{T}_1^a \mathbf{T}_2^b \mathbf{T}_3^c \mathbf{T}_4^d \left[f^{ade} f^{cbe} L_{1234}^2 \left(L_{1423} L_{1342}^2 + L_{1423}^2 L_{1342} \right) + \text{cycl.} \right]$$

where $L_{ijkl} = \log(\rho_{ijkl})$.



A three-loop diagram for Δ

- In the high-energy limit one finds (with $L = \log|s/t|$)

$$\rho_{1234} \equiv \frac{(-s_{12})(-s_{34})}{(-s_{13})(-s_{24})} = \left(\frac{s}{-t}\right)^2 \mathrm{e}^{-2\mathrm{i}\pi}\,; \qquad L_{1234} = 2(L - \mathrm{i}\pi)\,;$$

$$\rho_{1342} \equiv \frac{(-s_{13})(-s_{24})}{(-s_{14})(-s_{23})} = \left(\frac{-t}{s+t}\right)^2; \qquad L_{1342} \simeq -2L\,;$$

$$\rho_{1423} \equiv \frac{(-s_{14})(-s_{23})}{(-s_{12})(-s_{34})} = \left(\frac{s+t}{s}\right)^2 \mathrm{e}^{2\mathrm{i}\pi}\,; \qquad L_{1423} \simeq 2\mathrm{i}\pi\,,$$

- Previously admissible corrections display superleading high-energy logarithms at three loops.

$$\Delta^{(212)}(\rho_{ijkl}, \alpha_s)) = \left(\frac{\alpha_s}{\pi}\right)^3 \mathbf{T}_1^a \mathbf{T}_2^b \mathbf{T}_3^c \mathbf{T}_4^d \, 32\, \mathrm{i}\, \pi \left[\left(-L^4 - \mathrm{i}\pi L^3 - \pi^2 L^2 - \mathrm{i}\pi^3 L \right) f^{ade} f^{cbe} + \dots \right]$$

- No known explicit example of admissible quadrupole correction survives. A complete proof is still lacking: linear combinations might restore the proper Regge behavior.

OUTLOOK

Summary

- After 7.5 10^2 years, soft and collinear singularities in gauge theories amplitudes are still a fertile field of study. A definitive solution may be at hand.
  - ✓ We are probing the all-order structure of the nonabelian exponent.
  - ✓ All-order results constrain, test and complement fixed-order calculations.
  - ✓ Understanding singularities has phenomenological applications through resummation.

- Factorization theorems ⇒ Evolution equations ⇒ Exponentiation.
  - ✓ Sudakov factorization ⇒ soft-gluon resummation.
  - ✓ Multiparton processes require anomalous dimension matrices.

- A simple dipole formula may encode all infrared singularites for any massless gauge theory, a natural generalization of the planar limit. The study of possible corrections to the dipole formula is under way.

- The high-energy limit of the dipole formula provides insights into Reggeization and beyond, at least for divergent contributions to the amplitude.

- Leading logarithmic Reggeization is proved for generic color representations exchanged in the t channel, and for any number of partons in Multi-Regge kinematics.

- Regge factorization generically breaks down at NNLL, with computable corrections.

- The high-energy limit further constrains quadrupole corrections to the dipole formula: no known examples survive.

$$\boxed{\Gamma_{\rm dip}\left(\frac{p_i}{\mu}, \alpha_s(\mu^2)\right) = -\frac{1}{4}\,\widehat{\gamma}_K\left(\alpha_s(\mu^2)\right)\sum_{j\neq i}\ln\left(\frac{-2\,p_i\cdot p_j}{\mu^2}\right)\mathbf{T}_i\cdot\mathbf{T}_j + \sum_{i=1}^{n}\gamma_{J_i}\left(\alpha_s(\mu^2)\right)\,.}$$

INFRARED

CATASTROPHE

THANK YOU!