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High-energy resummation in Higgs production at the next-to-leading order

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We present the full next-to-leading order (NLO) result for the impact factor of a forward Higgs boson, obtained in the infinite-top-mass limit, both in the momentum representation and as superposition of the eigenfunctions of the leading-order (LO) BFKL kernel.

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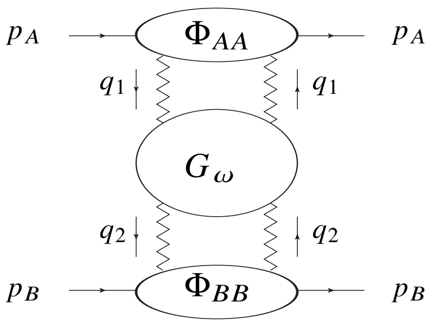

Figure 1: Schematic representation of the factorized amplitude.

1. Introduction

Precision physics in the Higgs sector has been one of the main challenges in recent years. The pure fixed-order calculations entering the *collinear factorization* framework, which have been pushed up to the N3LO, are not able to describe the entire kinematic spectrum. In particular conditions, they must be necessarily supplemented by all-order *resummations*; for instance, in the so called *Regge* kinematic region, large energy-type logarithms spoil the perturbative behavior of the series and must be resummed to all orders. This resummation is, for instance, necessary to describe the inclusive hadroproduction of a forward Higgs in the limit of small Bjorken x, as well as to study inclusive forward emissions of a Higgs boson in association with a backward identified object. In Refs. [1, 2], pioneering studies were performed on the Higgs production in mini-jet events within the leading-logarithmic approximation (LLA). Phenomenological LLA analyses on the Higgs plus jet(s) production were performed within partial next-to-leading logarithmic approximation (NLLA) in Refs. [3, 4]. High-energy effects from BFKL and Sudakov contributions were combined together to describe cross sections for the Higgs-plus-jet hadroproduction in almost back-to-back configurations [5]. Azimuthal correlations between a single-charmed hadron emitted in ultra-forward directions of rapidity and a Higgs boson were investigated within partial NLA in Ref. [6].

Nevertheless, a complete resummation for these processes at full NLLA can be achieved through the Balitsky-Fadin-Kuraev-Lipatov (BFKL) approach [7–10] (see Refs. [11–23] for recent applications), but it requires the knowledge of the next-to-leading order Higgs impact factor. We present the full NLO result for the impact factor of a forward Higgs boson, obtained in the infinite-top-mass limit, both in the momentum representation and as superposition of the eigenfunctions of the LO BFKL kernel.

2. BFKL approach

The BFKL equation is an integral equation that determines the behaviour at high energy $\sqrt{s}$ of the perturbative QCD amplitudes. It was derived in the LLA, which means collection of all terms of the type $\alpha_s^n \ln^n s$. This approximation leads to an increase of cross sections of the type

$$\sigma_{\rm tot}^{\rm LLA} = \frac{s^{\omega_0}}{\sqrt{\ln s}} \, , \tag{1}$$

where $\omega_0 = \frac{g^2 C_A \ln 2}{\pi^2}$ is the rightmost singularity in the complex momentum plane of the t-channel partial wave with vacuum quantum numbers (Pomerančuk singularity).

In the BFKL approach, using the s-channel unitarity relation, the imaginary part of a generic elastic scattering amplitude, $\mathcal{A}_{AB}^{AB}$, can be presented as[1]

$$\Im_s \mathcal{A}_{AB}^{AB} = s \int_{\delta - i\infty}^{\delta + i\infty} \frac{d\omega}{2\pi i} \int \frac{d^{D-2}\vec{q}_1 d^{D-2}\vec{q}_2}{(2\pi)^{D-2}} \left(\frac{s}{s_0} \right)^{\omega} \frac{\Phi_{AA}(\vec{q}_1)}{\vec{q}_1^{\,2}} G_\omega(\vec{q}_1, \vec{q}_2) \frac{\Phi_{BB}(-\vec{q}_2)}{\vec{q}_2^{\,2}}, \tag{2}$$

where $G_\omega(\vec{q}_1, \vec{q}_2)$ is the Mellin transform of the BFKL Green's function, which satisfies the BFKL equation, $\Phi_{AA}(\vec{q}_1)$ and $\Phi_{BB}(\vec{q}_2)$ are the so-called *impact factors*, δ is a real number which lies to the right of the right-most singularity of $G_\omega(\vec{q}_1, \vec{q}_2)$ and s_0 is scale introduced when performing the Mellin transform.

A schematic representation of the factorization formula (2) is given in Fig. 1. Although the s-behavior is governed by the Green's function, the impact factors are necessary ingredients to construct the total amplitude. In this work, we will be interested in computing a NLO impact factor, for this reason, we report here the full next-to-leading definition:

$$\Phi_{AA}(\vec{q}_1; s_0) = \left(\frac{s_0}{\vec{q}_1^{\,2}} \right)^{\omega(-\vec{q}_1^{\,2})} \sum_{\{f\}} \int \theta(s_\Lambda - s_{AR}) \frac{d s_{AR}}{2\pi} \, d\rho_f \, \Gamma_{\{f\}A}^{c} \left(\Gamma_{\{f\}A}^{c'} \right)^* \langle cc' | \hat{\mathcal{P}}_0 | 0 \rangle$$
$$- \frac{1}{2} \int d^{D-2} q_2 \, \frac{\vec{q}_1^{\,2}}{\vec{q}_2^{\,2}} \, \Phi_{AA}^{(0)}(\vec{q}_2) \, \mathcal{K}_r^{(0)}(\vec{q}_2, \vec{q}_1) \, \ln \left(\frac{s_\Lambda^2}{s_0 (\vec{q}_2 - \vec{q}_1)^2} \right) , \tag{3}$$

where

$$\mathcal{K}_r^{(0)}(\vec{q}_2, \vec{q}_1) = \frac{2 g^2 C_A}{(2\pi)^{D-1}} \frac{1}{(\vec{q}_1 - \vec{q}_2)^2} \tag{4}$$

is the real part of the leading order BFKL kernel and

$$\omega^{(1)}(t) = \frac{g^2 t}{(2\pi)^{D-1}} \frac{N}{2} \int \frac{d^{D-2} k_\perp}{k_\perp^2 (q-k)_\perp^2} = - \frac{g^2 C_A \Gamma(1+\epsilon) (\vec{q}^{\,2})^{-\epsilon}}{(4\pi)^{2-\epsilon}} \frac{\Gamma^2(-\epsilon)}{\Gamma(-2\epsilon)} , \tag{5}$$

is the one-loop Regge trajectory with $t = q^2 = -\vec{q}^{\,2}$. The impact factor in Eq. (3) is obtained by squaring off-shell amplitudes[2], $\Gamma_{\{f\}A}^{c}$, summing over all intermediate states, $\{f\}$, and integrating over the phase-space of the intermediate particles, $d\rho_f$, and over the invariant mass s_{AR} (invariant mass of the initial particle-Reggeon system). The factor $\langle cc' | \hat{\mathcal{P}}_0 | 0 \rangle$ is necessary to project onto the color-singlet representation. It is important to mention that the first line in Eq. (3) contains *rapidity divergences* when the invariant mass s_{AR} goes to infinity. From a technical point of view, the appearance of these divergences is due to the separation between the multi-Regge (MRK) and the quasi multi-Regge (QMRK) kinematics in the NLLA formulation of the BFKL approach. These divergences are regularized by the θ-function in the first line of Eq. (3) and then cancel with the

[1]By making use of the optical theorem, this can be further related to total cross-sections.
[2]The t-channel Reggeon is off-shell.

second line in the same equation. This latter contribution comes exactly from the MRK contribution to the amplitude (2) in the NLLA.

3. Effective gluon–Higgs coupling and LO impact factor

The calculation of the impact factor can be greatly simplified in the infinite-top-mass approximation. In this limit, we can employ the effective Lagrangian

$$\mathcal{L}_{ggH} = -\frac{g_H}{4} F^a_{\mu\nu} F^{\mu\nu,a} H \; , \qquad g_H = \frac{\alpha_s}{3\pi v}\left(1 + \frac{11}{4}\frac{\alpha_s}{\pi}\right) + \mathcal{O}(\alpha_s^3) \; , \tag{6}$$

which couples the Higgs field to gluons directly, via the QCD field strength tensor $F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g f^{abc} A^b_\mu A^c_\nu$.

At LO, the process is initiated by a collinear gluon that couples with the t-channel Reggeon to produce the forward Higgs boson. The gluon-initiated impact factor can be computed by using the definition given in Eq. (3). Then, the proton-initiated impact factor is related to the gluon-initiated one by the factorization

$$\frac{d\Phi^H_{PP}}{dx_H d^2\vec{p}_H} = \int_{x_H}^1 \frac{dz_H}{z_H} f_g\left(\frac{x_H}{z_H}\right) \frac{d\Phi^H_{gg}}{dz_H d^2\vec{p}_H} \; ,$$

where $\vec{p}_H$ is the transverse momentum of the Higgs with respect to the collision axis, z_H (x_H) is the longitudinal momentum fraction of the gluon (proton), and f_g is the gluon parton distribution function (PDF). At LO, the forward-Higgs impact factor reads

$$\frac{d\Phi^{\{H\}(0)}_{PP}}{dx_H d^2\vec{p}_H} = \frac{g_H^2 \vec{q}^{\,2} f_g(x_H) \delta^{(2)}(\vec{q} - \vec{p}_H)}{8\sqrt{N^2-1}} \; , \tag{7}$$

where $\vec{q}$ is the transverse momentum exchanged in the t-channel. This result correctly reproduces the large m_t-limit of the fully top-mass dependent result of Refs. [1, 3].

In the next sections, we employ the effective Lagrangian in Eq. (6) to compute the NLO corrections to the Higgs impact factor in the infinite-top-mass limit.

4. Next-to-leading order result in the k_T-space

4.1 Real corrections

At the next-to-leading order, the partonic sub-process can be initiated by a quark or a gluon. The quark-initiated contributions reads

$$\frac{d\Phi^{\{Hq\}}_{qq}(z_H, \vec{p}_H, \vec{q})}{dz_H d^2\vec{p}_H} = \frac{\sqrt{C_A^2 - 1}}{16 C_A (2\pi)^{D-1}} \frac{g^2 g_H^2}{(\vec{r}^{\,2})^2} \left[\frac{4(1-z_H)\,(\vec{r}\cdot\vec{q}\,)^2 + z_H^2 \vec{q}^{\,2} \vec{r}^{\,2}}{z_H} \right] \; , \tag{8}$$

while, the gluon-initiated one reads[3]

$$\frac{d\Phi_{gg}^{\{Hg\}}(\vec{q}\,)}{dz_H d^2\vec{p}_H} = \frac{g^2 g_H^2 C_A}{8(2\pi)^{D-1}(1-\epsilon)\sqrt{N^2-1}} \frac{2\vec{q}^{\,2}}{\vec{r}^{\,2}}$$

$$\times \left[\frac{z_H}{1-z_H} + z_H(1-z_H) + 2(1-\epsilon)\frac{(1-z_H)}{z_H}\frac{(\vec{q}\cdot\vec{r})^2}{\vec{q}^{\,2}\vec{r}^{\,2}} \right] \theta\left(s_\Lambda - s_{gR}\right) + \text{finite terms}\,, \qquad (9)$$

where $\vec{r} = \vec{q} - \vec{p}_H$. In Eq. (9) we have shown only the phase-space singular part of the impact factor; the complete result can be found in Ref. [24]. There are three kinds of phase-space singularities in Eqs. (8) and (9):

- Rapidity divergences when $z_H \to 1$, present in the gluon-initiated contribution.
- Soft divergences when $\vec{r} \to 0$ and $z_H \to 1$, present in the gluon-initiated contribution.
- Collinear divergences when $\vec{r} \to 0$, present both in the gluon-initiated and in the quark-initiated contribution.

Results in Eqs. (8) and (9) agree with ones in Ref. [25], independently performed in the Lipatov effective-action framework.

4.2 Virtual corrections

Being impact factors obtained by squaring effective vertices, we must extract the 1-loop effective vertex for the production of a Higgs in gluon-Reggeon collisions. To this aim, we use a reference amplitude and compare it with the expected Regge form. We employ the amplitude for the diffusion of a gluon off a quark to produce a Higgs plus a quark, $\mathcal{A}_{gq\to Hq}^{(8,-)}$, with octet color state and negative signature in the t-channel. It should assume the following Reggeized form[4]

$$\mathcal{A}_{gq\to Hq}^{(8,-)} = \Gamma_{Hg}^{ac}\frac{s}{t}\left[\left(\frac{s}{-t}\right)^{\omega(t)} + \left(\frac{-s}{-t}\right)^{\omega(t)}\right]\Gamma_{qq}^{c} \approx \Gamma_{Hg}^{ac(0)}\frac{2s}{t}\Gamma_{qq}^{c(0)}$$

$$+\Gamma_{Hg}^{ac(0)}\frac{s}{t}\omega^{(1)}(t)\left[\ln\left(\frac{s}{-t}\right) + \ln\left(\frac{-s}{-t}\right)\right]\Gamma_{qq}^{c(0)} + \Gamma_{Hg}^{ac(0)}\frac{2s}{t}\Gamma_{qq}^{c(1)} + \Gamma_{Hg}^{ac(1)}\frac{2s}{t}\Gamma_{qq}^{c(0)}\,, \qquad (10)$$

where $\omega(t)$ is the Regge trajectory, Γ_{Hg}^{ac} is the LO gluon-Higgs-Reggeon effective vertex and lastly Γ_{qq}^{c} is the quark-quark-Reggeon effective vertex. Since the only unknown ingredient in the right-hand side of Eq. (10) is the one-loop correction to the gluon-Higgs-Reggeon vertex, if we compute the amplitude $\mathcal{A}_{gq\to Hq}^{(8,-)}$, we are immediately able to extract it. The effective vertex allows us to obtain the virtual contribution to the impact factor, which reads[5]

$$\frac{d\Phi_{gg}^{\{H\}(1)}}{dz_H d^2\vec{p}_H} = \frac{d\Phi_{gg}^{\{H\}(0)}}{dz_H d^2\vec{p}_H}\frac{\bar{\alpha}_s}{2\pi}\left(\frac{\vec{q}^{\,2}}{\mu^2}\right)^{-\epsilon}\left[-\frac{C_A}{\epsilon^2} + \frac{11C_A - 2n_f}{6\epsilon}\right.$$

$$\left. -\frac{C_A}{\epsilon}\ln\left(\frac{\vec{q}^{\,2}}{s_0}\right) - \frac{5n_f}{9} + C_A\left(2\,\Re\left(\text{Li}_2\left(1+\frac{m_H^2}{\vec{q}^{\,2}}\right)\right) + \frac{\pi^2}{3} + \frac{67}{18}\right) + 11\right]\,. \qquad (11)$$

[3]This contribution does not contain the second line in Eq. (3).
[4]The apexes (0) and (1) denote the Born and the 1-loop approximation, respectively.
[5]For further details about the computation, see [24, 26, 27].

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In the computation $\epsilon = \epsilon_{UV} = \epsilon_{IR}$ is set such that any scaleless Feynman integral does not contribute to virtual corrections. The result in Eq. (11) is compatible with the independent result of Ref. [28], performed within the Lipatov effective action framework[6].

5. Next-to-leading order result in the (n, ν)-space and cancellation of divergences

5.1 The (n, ν)-space result

The cancellation of divergences can be seen only performing the integration of transverse momenta in Eq. (3). Nevertheless, in order to avoid a complete convolution between two impact factors and the Green function in the k_T-space, which would be complicated and would not lead to a fully general result[7], one can move to the (n, ν)-space. We explain the procedure, in the LLA, for pedagogical purpose. The BFKL Green function in Eq. (3) can be represent trough a spectral representation onto the eigenfunctions of its LO kernel, as

$$G_{\omega}^{(0)}(\vec{q}_1, \vec{q}_2) = \sum_{n=-\infty}^{\infty} \int_{-\infty}^{+\infty} d\nu \frac{\phi_{\nu}^{n}(\vec{q}_1^{\,2}) \phi_{\nu}^{n*}(\vec{q}_2^{\,2})}{\omega - \frac{\alpha_s C_A}{\pi} \chi(n, \nu)} \;, \tag{12}$$

where $\phi_{\nu}^{n}(\vec{q}^{\,2})$ are the LO BFKL kernel eigenfunctions and $(\alpha_s C_A/\pi)\chi(n, \nu)$ the corresponding eigenvalues, with

$$\phi_{\nu}^{n}(\vec{q}^{\,2}) = \frac{1}{\pi\sqrt{2}} (\vec{q}^{\,2})^{i\nu - \frac{1}{2}} e^{in\phi} \;, \qquad \chi(n, \nu) = 2\psi(1) - \psi\left(\frac{n}{2} + \frac{1}{2} + i\nu\right) - \psi\left(\frac{n}{2} + \frac{1}{2} - i\nu\right) \;. \tag{13}$$

Now, each impact factor integrate separately with one eigenfunction and we can give the definition of the (n, ν)-space projected impact factor,

$$d\Phi_{AA}^{(0)}(n, \nu) \equiv \int \frac{d^{2-2\epsilon} q}{\pi\sqrt{2}} (\vec{q}^{\,2})^{i\nu - \frac{3}{2}} e^{in\phi} d\Phi_{AA}^{(0)}(\vec{q}\,) \;. \tag{14}$$

Then, the projected LO Higgs impact factor reads

$$\frac{d\Phi_{PP}^{\{H\}(0)}(x_H, \vec{p}_H, n, \nu)}{dx_H d^2\vec{p}_H} = \frac{g_H^2}{8(1-\epsilon)\sqrt{N^2-1}} \frac{(\vec{p}_H^{\,2})^{i\nu - \frac{1}{2}} e^{in\phi_H}}{\pi\sqrt{2}} f_g(x_H) \;. \tag{15}$$

At NLO, the integration in (14) convert phase-space singularities of reals corrections in ϵ-poles, allowing to observe the cancellation of divergences.

5.2 Cancellation of divergences

Individual partonic contributions to the impact factor are clearly divergent. In order to build a finite quantity, we have to show the explicit cancellation of these divergences. We start discussing the rapidity divergences typical of high-energy computations. These are generated by the separation between MRK and QMRK; they are only present in the gluon-initiated contribution and give a

[6]To compare the two results, a reference physical amplitude must be considered; it must be reconstructed in both approaches. This is because there is freedom in how individual impact factors are defined.

[7]It would depend on the target impact factor.

$\ln s_\Lambda$-term after integration over the phase space. For this reason, the $d\Phi_{PP}^{\{Hg\}}$ contribution must be combined with the second line in Eq. (3). Symbolically, we then have

$$d\tilde{\Phi}_{PP}^{\{Hg\}} = d\Phi_{PP}^{\{Hg\}} - d\Phi_{PP}^{\{H\}} \otimes \mathcal{K}_r^{(0)} \ln s_\Lambda \, , \tag{16}$$

where $d\tilde{\Phi}_{PP}^{\{Hg\}}$ is free from rapidity divergences. This cancellation leave us with a double ϵ-pole singularity, coming from the kinematical region where $z_H \sim 1$ and $\vec{r} \sim \vec{0}$, after the integration over phase-space.

Virtual corrections are affected by UV-divergences, which can be removed performing the renormalization of the strong coupling, *i.e.*

$$\alpha_s(\mu^2) = \alpha_s(\mu_R^2) \left[1 + \frac{\alpha_s(\mu_R^2)}{2\pi} \beta_0 \left(-\frac{1}{\epsilon} - \ln(4\pi e^{-\gamma_E}) + \ln\left(\frac{\mu_R^2}{\mu^2} \right) \right) \right] . \tag{17}$$

Soft divergences should cancel in the real plus virtual combination, but, some IR-singularity (of collinear nature), survive this cancellation. They are initial state divergences that, within this scheme[8], should be cancelled by renormalizing the gluon PDF, *i.e.*

$$f_g(x,\mu) = f_g(x,\mu_F) - \frac{\alpha_s(\mu_F)}{2\pi} \left(-\frac{1}{\epsilon} - \ln(4\pi e^{-\gamma_E}) + \ln\left(\frac{\mu_F^2}{\mu^2} \right) \right)$$
$$\times \int_x^1 \frac{dz}{z} \left[P_{gq}(z) \sum_{a=q\bar{q}} f_a\left(\frac{x}{z}, \mu_F \right) + P_{gg}(z) f_g\left(\frac{x}{z}, \mu_F \right) \right] . \tag{18}$$

The cancellation of divergences takes place and the complete finite result can be expressed in terms of integrals of hypergeometric functions. We refer to the original work [24] for the complete result.

6. Conclusions and outlook

We calculated the full NLO correction to the impact factor for the production of a Higgs boson emitted by a proton in the forward rapidity region. Its analytic expression was obtained both in the momentum and in the Mellin representations. The latter is particularly relevant to clearly observe a complete cancellation of NLO singularities, and it is useful for future numerical studies [29].

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