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Forward Higgs production at NLO using Lipatov's high energy effective action

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Abstract

We use Lipatov's high energy effective action to determine the next-to-leading order corrections to Higgs production in the forward region. As a new element we provide a proper definition of the desired next-to-leading order coefficient within the high energy effective action framework, extending a previously proposed prescription. We further propose a subtraction mechanism to achieve for this coefficient a stable cancellation of real and virtual infra-red singularities in the presence of external off-shell legs and discuss aspects related to choice of a reference scale for high energy resummation.

Contents

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1 Introduction

In this contribution we present our result on the next-to-leading order corrections to the coefficient for the production of a Higgs boson using the framework of high energy factorization.

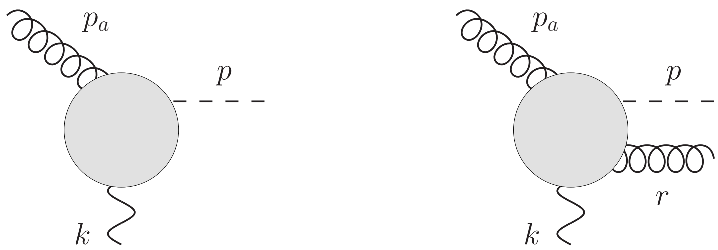

Figure 1: Kinematics for the leading order and virtual corrections (left); real next-to-leading order corrections (right).For a NLO quark final state, the intial and final gluon with momenta p_a and r respectively are to be replaced bt quark lines. The lines p_a and k are in-going, while p and r are out-going; figure taken from [7].

In particular our result is suitable to describe production of a Higgs boson in the forward direction of one of the scattering protons. The results presented here have been derived in [7] and make use of high energy effective action proposed by Lipatov [1] as well as the calculational framework developed in [2–6]. While our result has a direct application for the description of high energy resummed Higgs-jet configurations, as studied in [8], it is further of importance for the further development of the study of transverse momentum dependent (TMD) splitting kernels [9–11]and the TMD gluon distribution at low x [12], for which the Higgs boson provides a suitable color singlet state.

2 Conventions

To settle our conventions, we consider a process in which two hadrons A and B with light-like momenta $p_{A,B} = p^{\pm}_{A,B} n^{\mp}/2$ are scattering. The hadron momenta serve then to define light-cone directions $n^{\pm}$; the center of mass energy is within our conventions obtained as $\sqrt{s} = \sqrt{p_A^+ p_B^-}$. In the following we will present our results for inclusive production of an on-shell Higgs boson in the fragmentation region of hadron A; partons of hadron A are therefore characterized by relatively large values of hadron momentum fractions $x_A \sim 0.1$, while partons of hadron B have small momentum fractions $x_B < 10^{-2}$, which implies the necessity to resum logarithmically enhanced terms $\alpha_s \ln(1/x_B)$, where α_S denotes the strong coupling constant. The four momentum of the Higgs boson p and its rapidity η_H are parameterized as

$$p = x_H p_A + \frac{M_H^2 + \boldsymbol{p}^2}{x_H s} p_B + p_T, \qquad \eta_H = \ln \frac{x_H p_A^+}{\sqrt{M_H^2 + \boldsymbol{p}^2}}, \tag{1}$$

where $p_T^2 = -\boldsymbol{p}^2$ is the Higgs transverse momentum squared and x_H the proton momentum fraction, see also Fig. 1. To describe the coupling of the Higgs boson to the gluonic field, we make use of the heavy top limit which leads to the following effective Lagrangian, which describes the coupling of a Higgs boson to the gluonic field [13, 14],

$$\mathcal{L}_{\text{eff}} = -\frac{1}{4} g_H H F^a_{\mu\nu} F_a^{\mu\nu}. \tag{2}$$

Here H denotes the scalar (Higgs) field and g_H the effective Higgs-gluon coupling, which at reads at 1-loop as [15, 16]

$$g_H = -\frac{\alpha_s}{3\pi v}\left(1 + \frac{\alpha_s}{4\pi} 11\right) + \mathcal{O}(\alpha_s^3). \tag{3}$$

Since the top quarks is no longer a dynamical degree of freedom, the strong coupling α_s is evaluated for $n_f = 5$ flavors; $v^2 = 1/(\sqrt{2}G_F)$ and G_F denotes the Fermi constant. In the following we present our results which have been obtained under the assumption that multiple (reggeized) gluon exchange is still subleading. We have

$$\frac{d^3\sigma}{d^2\boldsymbol{p}dx_H} = \int_{x_H}^1 \frac{dz}{z} \sum_{a=q,g} f_a\left(\frac{x_H}{z}, \mu_F^2\right) \int \frac{d^2\boldsymbol{k}}{\pi} \frac{d\hat{C}_{ag^*\to H}(\mu_F^2, \eta_a; z, \boldsymbol{k})}{d^2\boldsymbol{p}dx_H} \mathcal{G}^{(B)}(\eta_a, \boldsymbol{k}), \tag{4}$$

where $\mathcal{G}^{(B)}(\eta_a, \boldsymbol{k})$ denotes the unintegrated gluon distribution of hadron B. It parameterizes non-perturbative input of hadron B and is subject to BFKL evolution; η_a is a factorization parameter associated with the highest gluon rapidity absorbed into $\mathcal{G}^{(B)}(\eta_a, \boldsymbol{k})$. For our convention for momenta see Fig. 1. We further use $0 < z < 1$ to parameterize the momentum fraction of the intial parton, carried on by the Higgs particle.

3 Results

We present our result for the differential cross section,

$$\frac{d^3\sigma}{dx_H d^2\boldsymbol{p}} = \int \frac{d^2\boldsymbol{r}}{\pi} \int_{x_H}^1 \frac{dz}{z} \sum_{a=q,g} f_a\left(\frac{x_H}{z}, \mu_F^2\right) \frac{d^3\hat{C}^{NLO}_{ag^*\to H}}{dx_H d^2\boldsymbol{p}} \mathcal{G}\left(\eta_a, (\boldsymbol{r}+\boldsymbol{p})^2\right) \tag{5}$$

where $\mathcal{G}(\eta_a, \boldsymbol{k})$ is the transverse momentum dependent unintegrated gluon density with $\boldsymbol{k}$ the transverse momentum and η_a the evolution parameter. The results for the coefficient were originally published in [7]. Here we reproduce these results, correcting two minor typing errors:

$$\frac{d^3\hat{C}^{NLO}_{ag^*\to H}}{dx_H d^2\boldsymbol{p}} = \sigma_0\left(\frac{d^3\hat{C}^{(0)}_{ag^*\to H}}{dx_H d^2\boldsymbol{p}} + \frac{\alpha_s}{2\pi}\frac{d^3\hat{C}^{(1)}_{g^*\to H}}{dx_H d^2\boldsymbol{p}} + \ldots\right), \qquad a = q, g \tag{6}$$

with

$$\frac{d^3\hat{C}^{(0)}_{gg^*\to H}}{dx_H d^2\boldsymbol{p}} = \delta^{(2)}(\boldsymbol{r})\delta(1-z), \qquad \frac{d^3\hat{C}^{(0)}_{qg^*\to H}}{dx_H d^2\boldsymbol{p}} = 0, \tag{7}$$

and

$$\begin{aligned}
&\frac{d^3\hat{C}^{(1)}_{gg^*\to H}(z, \boldsymbol{r}, \boldsymbol{p}; \eta_a, \mu_F, \mu)}{dx_H d^2\boldsymbol{p}} = \\
&= \delta(1-z)\Bigg\{2C_A\Bigg[\left(\ln\frac{(1-x_H)p_A^+}{\sqrt{\boldsymbol{r}^2}} - \eta_a\right)\frac{1}{\boldsymbol{r}^2}\Bigg]_+ + \delta^{(2)}(\boldsymbol{r})\Bigg[11 - \frac{\beta_0}{2}\ln\frac{\boldsymbol{p}^2}{\mu^2} - \frac{5n_f}{9} \\
&+ C_A\left(\frac{67}{18} + 2\mathrm{Li}_2\left(1 - \frac{M_H^2}{\boldsymbol{p}^2}\right) - \frac{\pi^2}{12}\right)\Bigg]\Bigg\} + \frac{H^{fin.}_{ggH}(z, \boldsymbol{p}, \boldsymbol{r})}{\boldsymbol{k}^2} + \left[\frac{\hat{P}^r_{gg}(z, \boldsymbol{p}, \boldsymbol{r})}{\boldsymbol{r}^2}\right]_+ \\
&+ \delta^{(2)}(\boldsymbol{r})\left[C_A\frac{1-z}{z}(4\ln 2 - 2) - \ln\frac{\mu_F^2}{\mu^2}P_{gg}(z)\right]
\end{aligned} \tag{8}$$

while

$$\int \frac{d^2\boldsymbol{r}}{\pi}\left[\frac{\kappa(\boldsymbol{r})}{\boldsymbol{r}^2}\right]_+ G((\boldsymbol{p}+\boldsymbol{r})^2) \equiv \int \frac{d^2\boldsymbol{r}}{\pi}\frac{\kappa(\boldsymbol{r})}{\boldsymbol{r}^2}\left[G(\boldsymbol{p}+\boldsymbol{r})^2) - \frac{\boldsymbol{p}^2 G(\boldsymbol{p}^2)}{\boldsymbol{r}^2 + (\boldsymbol{p}+\boldsymbol{r})^2}\right], \tag{9}$$

describes a novel plus prescription for convolutions in transverse momentum space, as the occur within high energy factorization. Finally

$$\frac{d^3\hat{C}^{(1)}_{qg^*\to H}(z,\boldsymbol{r},\boldsymbol{p};\eta_a,\mu_F,\mu)}{dx_H d^2\boldsymbol{p}} = \delta^{(2)}(\boldsymbol{r})\Big[C_F\frac{1-z}{z}(4\ln 2-4)-\ln\frac{\mu_F^2}{\mu^2}P_{gq}(z)\Big] + \frac{H^{fin.}_{qqH}(z,\boldsymbol{p},\boldsymbol{r})}{\boldsymbol{k}^2} + \left[\frac{\hat{P}^r_{qg}(z,\boldsymbol{p},\boldsymbol{r})}{\boldsymbol{r}^2}\right]_+ . \tag{10}$$

Furthermore

$$\begin{aligned} H^{fin.}_{ggH}(z,\boldsymbol{p},\boldsymbol{r}) &= \frac{2C_A}{z(1-z)}\left\{\frac{(1-z)^2z^2M_H^4}{2}\left(\frac{1}{\boldsymbol{\Delta}^2+(1-z)M_H^2}+\frac{1}{\boldsymbol{p}^2+(1-z)M_H^2}\right)^2\right. \\ &\quad -\frac{2z^2(\boldsymbol{p}\cdot\boldsymbol{\Delta})^2}{(\boldsymbol{p}^2+(1-z)M_H^2)(\boldsymbol{\Delta}^2+(1-z)M_H^2)}-\frac{2z(1-z)^2M_H^2}{\boldsymbol{\Delta}^2+(1-z)M_H^2} \\ &\quad -\frac{2z(1-z)^2M_H^2}{\boldsymbol{p}^2+(1-z)M_H^2}-\frac{(1-z)zM_H^2(\boldsymbol{k}\cdot\boldsymbol{r})z^2-2z^3(\boldsymbol{\Delta}\cdot\boldsymbol{r})(\boldsymbol{\Delta}\cdot\boldsymbol{k})}{\boldsymbol{r}^2(\boldsymbol{\Delta}^2+(1-z)M_H^2)} \\ &\quad \left.+\frac{(1-z)zM_H^2(\boldsymbol{k}\cdot\boldsymbol{r})z^2-2z^3(\boldsymbol{p}\cdot\boldsymbol{r})(\boldsymbol{p}\cdot\boldsymbol{k})}{\boldsymbol{r}^2(\boldsymbol{p}^2+(1-z)M_H^2)}+2z^2\right\} \\ &\quad +4C_A\frac{\boldsymbol{k}^2}{\boldsymbol{r}^2}\left[\frac{(1-z)}{z}\frac{(\boldsymbol{k}\cdot\boldsymbol{r})^2}{\boldsymbol{k}^2\boldsymbol{r}^2}-\frac{(1-z)}{z}\frac{(\boldsymbol{p}\cdot\boldsymbol{r})^2}{\boldsymbol{p}^2\boldsymbol{r}^2}\right], \\ H^{fin.}_{qqH}(z,\boldsymbol{p},\boldsymbol{r}) &= 4C_F\frac{\boldsymbol{k}^2}{\boldsymbol{r}^2}\left[\frac{(1-z)}{z}\frac{(\boldsymbol{k}\cdot\boldsymbol{r})^2}{\boldsymbol{k}^2\boldsymbol{r}^2}-\frac{(1-z)}{z}\frac{(\boldsymbol{p}\cdot\boldsymbol{r})^2}{\boldsymbol{p}^2\boldsymbol{r}^2}\right], \end{aligned} \tag{11}$$

and

$$\boldsymbol{\Delta} = z\boldsymbol{r}-(1-z)\boldsymbol{p} \qquad\qquad \boldsymbol{k} = \boldsymbol{p}+\boldsymbol{r}. \tag{12}$$

4 Conclusions

In this contribution we presented the next-to-leading order coefficient for forward Higgs production in the infinite top mass limit. Future studies will address its numerical properties as well as questions related to the setting of the high energy factorization scale.

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