THRESHOLD RESUMMATION FOR HIGGS PRODUCTION IN EFFECTIVE FIELD THEORY

FENG YUAN

RIKEN/BNL Research Center, Building 510A, Brookhaven National Laboratory, Upton, NY 11973

The threshold resummation effects for the Standard Model Higgs boson production at hadron collider are studied in the effective field theory formalism. The approach is conceptually simple, independent of details of an effective field theory formulation, and valid to all orders in sub-leading logarithms.

In hadron colliders, the rates on Higgs boson and Drell-Yan pair production demand reliable pQCD calculations. When the final-state invariant mass of hadrons is small, a fixed-order pQCD calculation yields large threshold double logarithms in the coefficient functions $\alpha_s^k \left[\frac{\ln^{m-1}(1-z)}{(1-z)}\right]_+$ $(m \leq 2k)$, which must be resummed to all orders in α_s, where $1-z$ is the fraction of center-of-mass energy of the initial partons going into soft radiations. In moment space, these large logarithms appear in the form, $\alpha_s^k \ln^m \overline{N}$, where $\overline{N} = N\exp(\gamma_E)$ with N, the order of moment. In the past decade, a standard method based on pQCD factorization has been established to perform the resummation [1,2].

In this talk, we introduced an alternative, effective field theory (EFT) resummation of these large threshold logarithms[3]. It is motivated by the recent development of soft-collinear effective field theory [4,5] and its applications to threshold resummation [6,7]. The basic idea to do resummation in EFT is two-step matching. At the higher scale, e.g., the Higgs mass, we match the gluon current between the full QCD and the EFT, from which the matching coefficients and the anomalous dimension can be calculated order by order. At lower scale, e.g., $M_H/\bar{N}$, the cross section (normally in moment-space) is matched to a product of parton distributions in the EFT. The scales of these matchings are chosen in such a way that both coefficients are free of large logarithms. In the following, I will summarize the resummation result for the Standard Model Higgs production in the

EFT. For the detailed derivation and the applications to other processes please refer to [3].

After integrating out the heavy quark loop, the Higgs boson production can be described by an effective lagrangian[8], $L = -1/4 C_\phi(M_t, \mu_R)\ \phi\ G^{\mu\nu}G_{\mu\nu}(\mu_R)$, where ϕ is the scalar field, $G^{\mu\nu}$ is the gluon field strength, C_ϕ is the effective coupling [9]. At higher scale, we match the gluon current $G^{\mu\nu}G_{\mu\nu}$ between the full QCD and the EFT. We introduce $a_s = \alpha_s/4\pi$ as expansion. Expanding the coefficient function at $\mu = M_H$ as $C_g(1, \alpha_s(M_H)) = \sum_i a_s^i(M_H) C_g^{(i)}$,

$$
\begin{aligned}
C_g^{(1)} &= 7 C_A \zeta_2 \\
C_g^{(2)} &= C_A^2 \left(\frac{5105}{162} + \frac{335}{6}\zeta_2 - \frac{143}{9}\zeta_3 + \frac{125}{10}\zeta_2^2 \right) \\
&\quad + C_A n_f \left(-\frac{916}{81} - \frac{25}{3}\zeta_2 - \frac{46}{9}\zeta_3 \right) + C_F n_f \left(-\frac{67}{6} + 8\zeta_3 \right) . \qquad (1)
\end{aligned}
$$

The relevant anomalous dimension of the gluon current can be written as

$$
\gamma_{1,g}^{(i)} = A_g^{(i)} \ln\left(M_H^2/\mu^2\right) + B_{1,g}^{(i)} + 2i\beta_{i-1} , \qquad (2)
$$

where $B_{1,g}^{(i)} = -2B_{2,g}^{(i)} - f_g^{(i)}$, A_g is the cusp anomalous dimension of Wilson lines in adjoint representation, and has been calculated to three-loops recently [10]. $B_{2,g}$ is the coefficient of $\delta(1-x)$ term in the gluon splitting function. The QCD β-function is defined as $\beta(a_s) = -d\ln\alpha_s/d\ln\mu^2 = \beta_0 a_s + \beta_1 a_s^2 + ...$ The functions $f_g^{(i)}$ are universal in the sense that the corresponding quark expressions are obtained by replacing the overall factor of C_A by C_F. Since $A^{(i)}$, $B_{2,g}^{(i)}$, and $f_g^{(i)}$ are known to three loops [11,12], the anomalous dimension is now known to the same order.

At the lower scale, one must consider soft-gluon radiations from the initial gluon partons. In principle, one should formulate a soft-collinear effective theory to calculate these contributions, as was done in Ref. [6]. However, this is unnecessary in practice and the result can simply be obtained from a full QCD calculation at the appropriate kinematic limit[13,14]. Expanding the matching coefficient, we get

$$
\begin{aligned}
M_N^{(1)} &= 2 C_A \zeta_2 \\
M_N^{(2)} &= C_A^2 \left[\frac{2428}{81} + \frac{67}{9}\zeta_2 - \frac{22}{9}\zeta_3 - 10\zeta_2^2 \right] \\
&\quad + C_A N_F \left[-\frac{328}{81} - \frac{10}{9}\zeta_2 + \frac{4}{9}\zeta_3 \right] , \qquad (3)
\end{aligned}
$$

at scale $\mu_I = M_H/\overline{N}$.

In the above results, QCD factorization produces gluon distributions at scale $\mu_I = M_H/\overline{N}$. We can bring the distributions to an arbitrary scale μ_F using the standard DGLAP evolution. This introduces an evolution factor, $\exp\left(2\int_{\mu_F}^{\mu_I}\frac{d\mu}{\mu}\gamma_{2,g}^N\right)$, where the twist-two anomalous dimension $\gamma_{2,g}^N$ has the following large N behavior, $\gamma_{2,g}^N = -A_g \ln \overline{N}^2 + 2B_{2,g}$, where A_g and $B_{2,g}$ are the same as those in Eqs. (8) and (9), respectively. To simplify the result, the factorization scale μ_F is henceforth chosen to be M_H.

Putting all factors together, the cross section in the moment-space is [15]

$$\sigma_N = \sigma_0 \cdot G_N(M_H) \cdot g(M_H, N) g(M_H, N) \ , \tag{4}$$

where σ_0 is a reference cross section and

$$G_N(M_H) = F(\alpha_s(M_H)) e^{I(\lambda, \alpha_s(M_H))} \tag{5}$$

where $F = |C(\alpha_s(M_H))|^2 M(\alpha_s(M_H))$ depends only on $\alpha_s(M_H)$. $I = I_1 + I_2 + I_3$ is a function of $\lambda = \beta_0 \ln \overline{N} \alpha_s(M_H)$ and $\alpha_s(M_H)$ with all leading and sub-leading large logarithms resummed, where $I_1 = 2\int_{M_H}^{\mu_I}\frac{d\mu}{\mu}\tilde{\gamma}_{1,g}$ with $\tilde{\gamma}_{1,g} = \gamma_{1,g} - 2i\beta_{i-1}$, is the anomalous dimension for $C = C_\phi \times C_g$, $I_2 = 2\int_{M_H}^{\mu_I}\frac{d\mu}{\mu}\gamma_{2,g}$, and $I_3 = -2\int_{\mu_I}^{M_H}\frac{d\mu}{\mu}\triangle B_1$, and $\triangle B_1$ is defined as $\triangle B_1 = -\beta(a_s) d\ln M_N / d\ln \alpha_s$.

The above result can be related to the conventional expression if one writes $I = I_\triangle + \ln \Delta C$, where

$$\begin{aligned} I_\triangle = \int_0^1 dz \frac{z^{N-1}-1}{1-z}\Bigg[2\int_{M_H^2}^{(1-z)^2 M_H^2} \frac{d\mu^2}{\mu^2} A_g\left(\alpha_s(\mu^2)\right) \\ + D_g(\alpha_s((1-z)^2 M_H^2))\Bigg] \ , \end{aligned} \tag{6}$$

and ΔC is just a function of $\alpha_s(M_H)$, serving to cancel the non-logarithmic terms in $I_\triangle$. Using similar methods as the ones in Ref. [15], it is a matter of some technical steps to get,

$$\begin{aligned} D_g(\mu^2) &= 2(B_{1,g} + \Delta B_1 + 2B_{2,g}) - \partial_{\alpha_s}\Gamma_2(\partial_{\alpha_s})\left[4A_g(\alpha_s) - \partial_{\alpha_s} D_g(\alpha_s)\right] \\ \Delta C &= \Gamma_2(\partial_{\alpha_s})\left[4A_g(\alpha_s) - \partial_{\alpha_s} D_g(\alpha_s)\right] \ , \end{aligned} \tag{7}$$

where $\Gamma_2(\epsilon) = 1/\epsilon^2[1 - e^{-\gamma_E\epsilon}\Gamma(1-\epsilon)] = -\zeta_2/2 - \zeta_3\epsilon/3 + ...$ and $\partial_{\alpha_s} = 2\beta(\alpha_s)\alpha_s\partial/\partial\alpha_s$. The above equations are our main result connecting the EFT resummation to the coventional approach, valid to all orders in leading and sub-leading logarithms. The $D_g^{(i)}$ coefficients can be solved iteratively from Eq. (7),

$$
\begin{aligned}
D_g^{(1)} &= 0 \\
D_g^{(2)} &= -2f_g^{(2)} + 4\beta_0\zeta_2 A_g^{(1)} - 2\beta_0 M_N^{(1)} \\
D_g^{(3)} &= -2f_g^{(3)} + 4\zeta_2\beta_1 A_g^{(1)} + 8\zeta_2\beta_0 A_g^{(2)} + \frac{32}{3}\zeta_3\beta_0^2 A_g^{(1)} \\
&\quad -2\beta_1 M_N^{(1)} - 2\beta_0\left[2M_N^{(2)} - \left(M_N^{(1)}\right)^2\right] \ ,
\end{aligned} \tag{8}
$$

which reproduces the recent calculations [16] in conventional approach.

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