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Higgs production in bottom-quark fusion in a matched scheme

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ABSTRACT

We compute the total cross-section for Higgs boson production in bottom-quark fusion using the so-called FONLL method for the matching of a scheme in which the b-quark is treated as a massless parton to that in which it is treated as a massive final-state particle. We discuss the general framework for the application of the FONLL method to this process, and then we present explicit expressions for the case in which the next-to-next-to-leading-log five-flavor scheme result is combined with the leading-order $\mathcal{O}(\alpha_s^2)$ four-flavor scheme computation. We compare our results in this case to the four- and five-flavor scheme computations, and to the so-called Santander matching.

In perturbative QCD, processes involving bottom quarks can be computed within different factorization schemes. One possibility is to use a five-flavor, or massless, scheme, in which the b-quark is treated as a massless parton. In this scheme, collinear logarithms of μ_F^2/m_b^2 (with μ_F the factorization scale) are resummed through QCD evolution equations, but corrections suppressed by powers of m_b^2/μ_F^2 are neglected. Alternatively, one may use a four-flavor, massive, or decoupling scheme, in which the b-quark is treated as a massive particle, which decouples from evolution equations and the running of α_s, but full dependence on m_b is retained. Generally, of course, results in the two schemes may differ by a large amount: indeed, the leading-order predictions for Higgs boson in bottom-quark fusion [1–4] may differ by up to one order of magnitude [5], though the disagreement is reduced if the factorization and renormalization scales are chosen to be smaller than m_H (which may well [6–10] be more appropriate) and higher perturbative orders are included.

The five-flavor scheme is more accurate for scales $\mu^2 \gg m_b^2$, while the four-flavor scheme is more accurate close to threshold, though of course if the four-flavor computation is performed to high enough order in perturbation theory it will reproduce the five-flavor scheme result (the converse is not true, because mass corrections are not included in the five-flavor scheme at any perturbative order). It is therefore advantageous to combine the two computations into one which is accurate at all scales. A phenomenological way of doing so, the so-called Santander matching, has been proposed in Ref. [11]: it consists of simply interpolating between the four- and five-flavor scheme results by means of a weighted average, such that in the two limits $\mu/m_b \gg 1$ or $\mu/m_b \sim 1$ the massless or massive results are respectively reproduced.

However, a more systematic approach which preserves the perturbative accuracy of both computations may be desirable. One such approach, the FONLL method, was proposed in Ref. [12] in the context of hadro-production of heavy quarks, and extended to deep-inelastic scattering in Ref. [13]. The basic idea of this method is to expand out the five-flavor-scheme computation in powers of the strong coupling α_s, and replace a finite number of terms with their massive-scheme counterparts. The result then retains the accuracy of both ingredients: at the massive level, the fixed-order accuracy corresponding to the number of massive orders which have been included (FO, or fixed order), and at the massless level, the logarithmic accuracy of the starting five-flavor scheme computation (NLL, or generally subleading logarithmic[1]).

It is the purpose of this paper to present the application of the FONLL scheme to Higgs production in bottom-quark fusion, focusing for definiteness on the total cross-section. In the rest of this paper we will follow the notation and conventions of Ref. [13].

[1] We will consistently use the notation N^{k}LL to refer to the resummation of collinear logs of the heavy quark mass, i.e. by LL we mean a computation in which $\left(\alpha_s \ln \frac{m_b^2}{\mu^2}\right)$ is treated as order one (α_s^0).

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The total cross-section σ in the five-flavor scheme has the form

$$\sigma^{(5)} = \iint dx_1 dx_2 \sum_{ij} f_i^{(5)}(x_1,\mu^2) f_j^{(5)}(x_2,\mu^2) \times \hat{\sigma}_{ij}^{(5)}\left(x_1,x_2,\alpha_s^{(5)}(\mu^2)\right), \quad (1)$$

where the sum runs over the 10 quarks and antiquarks and the gluon, and the b quark and antiquark are treated as the other partons, which in particular contribute to the running of $\alpha_s^{(5)}$. For simplicity we omit the dependence of the hard cross-section on the renormalization and factorization scales, which henceforth we will assume to be chosen equal to $\mu_R = \mu_F = \mu$, unless otherwise stated.

In the four-flavor scheme it has the form

$$\sigma^{(4)} = \iint dx_1 dx_2 \sum_{ij} f_i^{(4)}(x_1,\mu^2) f_j^{(4)}(x_2,\mu^2) \times \hat{\sigma}_{ij}^{(4)}\left(x_1,x_2,\frac{\mu^2}{m_b^2},\alpha_s^{(4)}(\mu^2)\right), \quad (2)$$

where now the sum only runs over the four lightest quarks and antiquarks and the gluon, the b-quark decouples from the running of $\alpha_s^{(4)}$ and the DGLAP evolution equations satisfied by $f_i^{(4)}(x_1,\mu^2)$, but full m_b dependence of the partonic cross-section $\hat{\sigma}_{ij}^{(4)}$ is retained.

In order to carry out the FONLL procedure, we need to express the four-flavor scheme cross-section, Eq. (2), in terms of $\alpha_s^{(5)}$ and $f_i^{(5)}$, so that their perturbative expansions can be compared directly. The coupling constant and the PDFs are related in the two schemes by equations of the form

$$\alpha_s^{(5)}(\mu^2) = \alpha_s^{(4)}(\mu^2) + \sum_{i=2}^{\infty} c_i(L) \times \left(\alpha_s^{(4)}(m_b^2)\right)^i, \quad (3)$$

$$f_i^{(5)}(x,\mu^2) = \int_x^1 \frac{dy}{y} \sum_j K_{ij}\left(y, L, \alpha_s^{(4)}(\mu^2)\right) f_j^{(4)}\left(\frac{x}{y},\mu^2\right), \quad (4)$$

where

$$L \equiv \ln \mu^2/m_b^2 \quad (5)$$

and the sum runs over the eight lightest flavors, antiflavors, and the gluon, while the index i takes value over all ten quarks and antiquarks and the gluon. The coefficients $c_i(L)$ are polynomials in L, and the functions K_{ij} can be expressed as an expansion in powers of α_s, with coefficients that are polynomials in L.

The first nine equations (4) relate the eight lightest quarks and the gluon in the two schemes and can be inverted to express the four-flavor-scheme PDFs in terms of the five-flavor-scheme ones. The last two equations, assuming that the bottom quark is generated by radiation from the gluon (i.e. no “intrinsic” [14] bottom component) express the bottom and anti-bottom PDFs in terms of the other ones. In particular, this assumption implies that the b quark and antiquark PDFs are equal to each other, $f_b^{(5)} = f_{\bar{b}}^{(5)}$. Inverting Eqs. (3)–(4) and substituting in Eq. (2) one can obtain an expression of $\sigma^{(4)}$ in terms of $\alpha_s^{(5)}$ and $f_i^{(5)}$:

$$\sigma^{(4)} = \iint dx_1 dx_2 \sum_{ij=q,g} f_i^{(5)}(x_1,\mu^2) f_j^{(5)}(x_2,\mu^2) \times B_{ij}^{(4)}\left(x_1,x_2,\frac{\mu^2}{m_b^2},\alpha_s^{(5)}(\mu^2)\right), \quad (6)$$

where the coefficient functions B_{ij} are such that substituting the matching relations Eqs. (3)–(4) in Eq. (6) the original expression Eq. (2) is recovered. Note that in the course of the procedure of expressing $\sigma^{(4)}$ in terms of $\alpha_s^{(5)}$ and $f_i^{(5)}$, subleading terms are introduced, because (3)–(4) are only inverted to finite perturbative accuracy. It follows that the expressions Eq. (2) and Eq. (6) of $\sigma^{(4)}$ actually differ by subleading terms. Henceforth, for $\sigma^{(4)}$ we will use the expression Eq. (6), and avoid any further reference to $\alpha_s^{(4)}$ and $f_i^{(4)}$; therefore, from now on α_s and f_i will denote the five-flavor scheme expressions.

In order to match the two expressions for σ in the five-flavor scheme, Eq. (1), and in the four-flavor scheme, Eq. (6), we now work out their perturbative expansion. Using DGLAP evolution, the b-PDF, $f_b^{(5)}(\mu^2)$, can be determined in terms of the gluon and the light-quark parton distributions $f_i^{(5)}$ at the scale μ^2 convoluted with coefficient functions expressed as a power series in $\alpha_s^{(5)}$, with coefficients that are polynomials in L. The five-flavor-scheme expression Eq. (1) may thus be written entirely in terms of light-quark and gluon PDFs:

$$\sigma^{(5)} = \iint dx_1 dx_2 \sum_{ij=q,g} f_i^{(5)}(x_1,\mu^2) f_j^{(5)}(x_2,\mu^2) \times A_{ij}^{(5)}\left(x_1,x_2,L,\alpha_s^{(5)}(\mu^2)\right), \quad (7)$$

where the $A_{ij}^{(5)}$ coefficient functions are given by a perturbative expansion of the form

$$A_{ij}^{(5)}\left(x_1,x_2,L,\alpha_s^{(5)}(\mu^2)\right) = \sum_{p=0}^{N}\left(\alpha_s^{(5)}(\mu^2)\right)^p \sum_{k=0}^{\infty} A_{ij}^{(p),(k)}(x_1,x_2)\left(\alpha_s^{(5)}(\mu^2)L\right)^k, \quad (8)$$

with at leading order $N=0$, and at N^{m}LO order $N=m$.

On the other hand, the four-flavor-scheme expression Eq. (6), as mentioned, is also written in terms of the light-quark PDFs, with coefficient functions B_{ij} which can also be expanded in power of $\alpha_s^{(5)}$,

$$B_{ij}^{(4)}\left(x_1,x_2,\frac{\mu^2}{m_b^2},\alpha_s^{(5)}(\mu^2)\right) = \sum_{p=0}^{N}\left(\alpha_s^{(5)}(\mu^2)\right)^p B_{ij}^{(p)}\left(x_1,x_2,\frac{\mu^2}{m_b^2}\right), \quad (9)$$

where N is the order of the expansion needed to reach the desired accuracy. It follows that the sum of all contributions to the four-flavor-scheme expression Eq. (9) which do not vanish when $\mu^2 \gg m_b^2$ must also be present in the five-flavor-scheme result.

These contributions $B_{ij}^{(0),(p)}$ provide the massless limit of $B_{ij}^{(p)}$, in the sense that

$$\lim_{m_b\to 0}\left[B_{ij}^{(p)}\left(x_1,x_2,\frac{\mu^2}{m_b^2}\right) - B_{ij}^{(0),(p)}\left(x_1,x_2,\frac{\mu^2}{m_b^2}\right)\right] = 0. \quad (10)$$

In other words, $B_{ij}^{(0),(p)}$ is obtained from $B_{ij}^{(p)}$ by retaining all logarithms and constant terms and dropping all terms suppressed by powers of m_b/μ. Given that these terms are also present in the five-flavor-scheme calculation, we can also write

$$B_{ij}^{(0),(p)}\left(x_1,x_2,\frac{\mu^2}{m_b^2}\right) = \sum_{k=0}^{p} A_{ij}^{(p-k),(k)}(x_1,x_2)\,L^k \quad (11)$$