The Heavy Quark Parton Oxymoron
– A mini-review of Heavy Quark Production theory in PQCD†

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Abstract. Conventional perturbative QCD calculations on the production of a heavy quark “H” consist of two contrasting approaches: the usual QCD parton formalism uses the zero-mass approximation \(m_H = 0\) once above threshold, and treats \(H\) just like the other light partons; on the other hand, most recent “NLO” heavy quark calculations treat \(m_H\) as a large parameter and always consider \(H\) as a heavy particle, never as a parton, irrespective of the energy scale of the physical process. By their very nature, both these approaches are limited in their regions of applicability. This dichotomy can be resolved in a unified general-mass variable-flavor-number scheme, which retains the \(m_H\) dependence at all energies, and which naturally reduces to the two conventional approaches in their respective region of validity. Recent applications to lepto- and hadro-production of heavy quarks are briefly summarized.

1 Introduction

The production of heavy quarks in photo-, lepto-, and hadro-production processes has become an increasingly important subject of study both theoretically and experimentally. For a comprehensive review and references, see Ref. [1]. The theory of heavy quark production in perturbative Quantum Chromodynamics (PQCD) is considerably more subtle than that of light parton (jet) production because of the additional scale introduced by the quark mass. Let us consider the production of a generic heavy quark, denoted by \(H\), with non-zero mass \(m_H\), in high energy interactions. For definiteness and simplicity, unless otherwise stated, we shall use deep inelastic lepton-hadron scattering as the talking example. A reasonable criterion for a quark to be called “heavy” is \(m_H \gg \Lambda_{QCD}\), so that perturbative QCD is applicable at the scale \(m_H\). Thus, conventionally, \(\{c, b, t\}\) quarks are regarded as heavy.

The relevant energy scales of this problem are: (i) a typical small scale such as \(\Lambda_{QCD}\) or masses of light mesons, nucleons, ...; (ii) the highest energy scale \(E\) or \(\sqrt{s}\); (iii) a typical
large scale in the physical process, such as $p_t$ of the heavy quark (or the associated heavy flavor hadron), $Q$ of deep inelastic scattering or Drell-Yan processes, or some large mass (such as $m_W, m_Z, m_{Higgs}, m_{SUSY}$) – to be denoted henceforth collectively as $Q$; and (iv) the heavy quark mass $m_H$. By definition, $m_H \gg \Lambda_{QCD}$; and we need $\sqrt{s}$ to be fairly large compared to $m_H$ for the production cross-section to be substantial. Thus the important ratio of scales remaining which determines the physics of the heavy quark production process is that between $m_H$ and $Q$.

We shall be mainly concerned with $c$ and $b$ quarks for which this ratio can vary over a wide range in practice.

2 Two Contrasting Conventional Approaches

The two conventional approaches to heavy quark production in PQCD can be summarized by the following contrasting master equations used in the calculation

\[
\text{ZM-VFN: } \sigma_{IA \rightarrow CX} = \sum_{a = \text{all active partons}} f_a^a(x_a, \mu) \otimes \hat{\sigma}_{la \rightarrow CX}(\hat{s}, Q, \mu) \bigg|_{\overline{\text{MS}}}^{m_a=0} (1)
\]

\[
\text{FFN: } \sigma_{IA \rightarrow HX} = \sum_{a = \text{light partons only}} f_a^a(x_a, \mu) \otimes \hat{\sigma}_{la \rightarrow HX}^{FFN}(\hat{s}, Q, m_H, \mu) (2)
\]

The **Zero-mass Variable-flavor-number (ZM-VFN) scheme** formula, Eq.1, is used routinely in most high energy calculations: in global analyses of parton distributions, from EHLQ [2] to MRS [3] and CTEQ [4], as well as in all analytic or Monte Carlo programs for generating SM and new physics cross-sections. In this equation, the parton label $a$ is summed over all possible active parton species; $\mu$ is the factorization and renormalization scale; and $\hat{\sigma}_{la \rightarrow CX}$ is the perturbatively calculable hard cross section involving partons only. “Active” partons include all quanta which can participate effectively in the dynamics at the relevant energy scale $\mu (\sim Q)$ [2,5,6], including charm and bottom quarks at current collider energies. Thus, the active flavor number $n_f$ depends on the energy scale (“resolving power”) of the problem; it is not fixed at any particular value. The hard cross-section $\hat{\sigma}_{la \rightarrow CX}(\hat{s}, Q, \mu)$ is calculated in the limit of zero mass for all the partons, and it is made infra-red safe by dimensional regularization in the $\overline{\text{MS}}$ scheme – hence the name **Zero-mass Variable-flavor-number (ZM-VFN) scheme**.

The advantage of the ZM-VFN scheme is that it is quite simple to implement. For the light partons $a = \{g, u, d, s\}$, $m_a \rightarrow 0$ is a valid approximation for all hard scale $Q$ (since, by definition, $Q \gg m_a$). But for a heavy quark $H$, it is a reasonable approximation only in the high energy regime $\mu \sim Q \gg m_H$; and it clearly becomes unreliable in the intermediate region $Q \sim \mathcal{O}(m_H)$.

1) In this talk, we shall not consider “small-x” problems associated with logarithms of the large ratio $\sqrt{s}/Q$, cf [12].

2) If a final state particle $C$ is observed, the factorization formula should also contain a fragmentation function $d^C(z, \mu)$. We leave out $d^C(z, \mu)$ here only for simplicity of discussion. All statements concerning the parton distributions also apply to the fragmentation functions, if present.
In contrast to the above, the fixed-flavor-number (FFN) scheme, Eq.2 has been used in most recent fixed-order perturbative calculations of heavy quark production [7–9]. In this scheme, by definition, only light partons (e.g. u, d, s and g for charm production) are included in the initial state: the number of parton flavors $n_f$ is kept at a fixed value regardless of the energy scales involved ($n_f = 3, 4$ for c, b production respectively). The main feature here is: $H$ is pictured as a heavy particle – much in the same way as W, Z, and other new heavy particles, and very different from the zero-mass light partons – hence the mass $m_H$ is kept exactly in the hard cross-section $\hat{\sigma}_{ab \rightarrow HX}^{FFN}(\hat{s}, Q, m_H, \mu)$. Typically, the perturbative $\hat{\sigma}_{ab \rightarrow HX}^{FFN}(\hat{s}, Q, m_H, \mu)$ will contain logarithm factors of the form $a_n^\alpha(\mu)\ln^{n-k-m}(Q/m_H)\ln^m(\hat{s}/Q^2)$. If $Q \sim m_H$ (and $x \sim Q^2/s$ is not too small), these factors are under control; and we have effectively a one large scale hard process. Hence, the FFN scheme is the natural scheme to use in the energy region $Q \sim m_H$ – this is precisely where the ZM VFN scheme is expected to be inappropriate.

From the heavy particle perspective, this approach also has the advantage of being conceptually simple, even if the NLO calculation requires considerable amount of work. However, the sharp distinction drawn between the $H$ quark and the other light quarks, say between $c$ and $s$, in this formalism appears quite unnatural as the hadron system is probed at the scale $\mu \sim Q$ available in current high energy processes. And it has been known since the next-to-leading order (NLO) calculations in the FFN scheme were completed [8,9] that, for both charm and bottom production, there are two disconcerting features about the results: (i) the NLO corrections turn out to be of the same numerical magnitude as (in fact, generally larger than) the leading order (LO) result; and (ii) the uncertainty of the theoretical calculation, as measured by the dependence of the calculated cross section on the unphysical scale parameter $\mu$, is as large in NLO as in LO – contrary to what is expected from a good perturbation expansion [10]. These features mean that the truncated perturbative series in this scheme has left out important physics effects. Experimentally, it is also known that the measured charm and bottom production cross sections do not agree with the NLO theoretical predictions, at least in the overall normalization, even when the scale $\mu$ is allowed to vary within a reasonable range [11].

This situation may not be all that surprising: for $c$ and $b$ quarks, the condition $Q \sim m_H$ is not well satisfied in most practical cases. In fact, current experimental ranges for lepto- and hadro-production of these heavy flavors mostly lie in a region between those appropriate for the ZM VFN ($Q \gg m_H$) and FFN ($Q \sim m_H$) schemes. We need a well-defined theory which applies over the full $Q$ range! Other possible sources for these problems are: (i) large corrections due to large logarithms of $(s/Q^2)$—the small-$x$ problem [12]; (ii) inadequate understanding of the hadronization of heavy quarks in comparing PQCD calculations with experiment; and (iii) existence of non-perturbative components of $H$ inside the nucleon which are, by definition, excluded by the FFN scheme. In this talk, we shall concentrate on physics issues pertaining to the changing role of the heavy quark $H$ over the full $Q$ range. It is particularly interesting because the interplay between the two independent scales $m_H$ and $Q$ embodies much interesting QCD physics which is amenable to precise treatment.
3 A Unified, General-mass, Parton Approach

When the energy scale becomes large, $Q/m_H \gg 1$, the FFN scheme becomes suspect because large logarithm factors of $\ln(Q/m_H)$ in the hard cross section $\hat{\sigma}_{ab \rightarrow HX}^{FFN}$ becomes increasingly singular, and higher-order terms containing higher powers of the same can no longer be omitted. In other words, the truncated perturbation series in this scheme can become rather unreliable as $Q$ becomes large. The clue for addressing this problem is already contained in Eq.1: large logarithms of the form $\alpha_s^n \ln^{-k}(Q/m_H)$ in $\hat{\sigma}_{ab \rightarrow HX}^{FFN}$ can be resummed to all orders in $\alpha_s$ to become the parton distribution $f_A^H(x, \mu)$ (evolved to $(k+1)$-loops). The $H$ parton should be included in the sum over parton flavors; it participates in the hard scattering on the same footing as the other partons. After removing these potentially dangerous logarithm terms, the remaining hard cross section $\hat{\sigma}_{ab \rightarrow HX}$ becomes infra-red safe as $Q/m_H \rightarrow \infty$. It is important to note, however, the resummation of large $\ln(Q/m_H)$ logarithms does not require taking the $m_H \rightarrow 0$ limit for the remaining (“mass-subtracted”) hard cross-section as is done in the conventional ZM VFN formulation, Eq.1. In fact, by retaining the $m_H$ dependence in the mass-subtracted (hence infra-red safe) $\hat{\sigma}_{ab \rightarrow HX}(\hat{s}, Q, m_H, \mu)$, one arrives at a consistent theory for heavy quark production which is valid over the entire energy range from $Q \lesssim m_H$ to $Q \gg m_H$:

$$\sigma_{lA \rightarrow CX}(s, Q, m_H) = \sum_{a = \text{all active partons}} f_a^a(x_a, \mu) \otimes \hat{\sigma}_{lA \rightarrow CX}(\hat{s}, Q, m_H, \mu)$$  (3)

A program to systematically implement this intuitive physical picture has been developed in a series of papers in [6,13–15]. The resulting formalism constitutes a natural generalization of the conventional zero-mass QCD parton formalism to correctly include general quark mass effects, hence will be called the general-mass variable-flavor-number (GM-VFN) scheme. (In some recent literature it has also been called the ACOT scheme, Ref. [15].)

More precisely, this formalism is based on a well defined renormalization scheme [16] which provides a natural transition from the threshold region $Q \sim O(m_H)$ to the high energy region $Q \gg m_H$; and the validity of the generalized factorization theorem can be established order-by-order in perturbation theory [17]. The key points are:

- the renormalization scheme is a composite of two simple schemes, natural for $Q \lesssim m_H$ and $Q \geq m_H$ respectively, with matching conditions that make the schemes equivalent in the domain of overlap $Q \sim m_H$ where they are equally valid for practical low order calculations [16];
- one scheme utilizes a subtraction procedure (BPHZ) which leads to manifest decoupling of the heavy particle in the region $Q \ll m_H$, thereby gives precise meaning to the FFN scheme (with no heavy quark partons);
- the other scheme is ordinary $\overline{MS}$ as regards the definition of the coupling $\alpha_s(\mu)$ and the parton densities $f_A^a(x, \mu)$, hence retains the normal ($m_H = 0$) evolution equations.

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3) Factorization of any applicable fragmentation functions is implicitly assumed.
for the latter in the region $Q \gtrsim m_H$ [6] – this comes about because the evolution kernels are anomalous dimensions which are derivatives of renormalization constants, and renormalization constants are mass-independent in the $\overline{\text{MS}}$ scheme;

- the factorization scheme is defined such that all infra-red safe $m_H$-dependent effects are preserved in the hard cross-sections, so that there is no loss of accuracy when $Q \sim m_H$ [15]. This is accomplished by defining $\hat{\sigma}_{l \to CX}(s, Q, m_H, \mu)$ as the full $\sigma_{l \to CX}(s, Q, m_H, \mu)$ with mass ($m_H$) singularities subtracted.

These features guarantee that predictions of this formalism: (a) coincide with those of the FFN scheme in its region of applicability, $Q \lesssim m_H$; (b) reduce to those of the conventional zero-mass parton model in the asymptotic energy regime $Q \gg m_H$; and (c) provide a good approximation to the physical cross-section over the entire energy range in between, since the remainder of the perturbation series contains no large logarithms of the kind $\log(Q/m_H)$.

4 How, and Why, does the ACOT scheme work?

How and why does this scheme work were described in some detail for lepto-production of heavy quarks in Ref. [15]. We recall the essential points here. Writing out the first two terms in the perturbative expansion of Eq.3, we have

$$\sigma_{lA \toHX}(s, Q, m_H) = f_A^H \otimes 0\hat{\sigma}_{lH \to lH} + f_A^0 \otimes 1\hat{\sigma}_{l g \to lH}$$(4)

where the superscript $n$ in $n\hat{\sigma}$ denotes the order in $\alpha_s$ of the hard cross-section $\hat{\sigma}$. The lowest order hard cross-section $0\hat{\sigma}_{lH \to lH}$ is identical to the Born expression $0\sigma_{lH \to lH}$ since the tree diagram does not need any subtraction. The order $\alpha_s$ hard cross-section is given by

$$1\hat{\sigma}_{l g \to lH} = 1\sigma_{l g \to lH} - 1 f_g^H \otimes 0\hat{\sigma}_{lH \to lH}$$  

$$1 f_g^H = \frac{\alpha_s(\mu)}{2\pi} \ln \frac{\mu^2}{m_H^2} P_{gH}$$ (5)

where $1 f_g^H$ is the perturbative parton distribution function of finding $H$ in $g$; $P_{gH}$ is the usual $g \to H$ splitting function; and $1\sigma_{l g \to lH}$ is the unsubtracted cross-section for the gluon fusion process $l g \to lH$. The subtraction term above can be formally derived by applying Eq.4 to the partonic cross-section $1\sigma_{l g \to lH}$ (with light-parton colinear singularities subtracted), and then solving for $1\hat{\sigma}_{l g \to lH}$ [15].

Substituting Eq.5 in Eq.4, the physical cross-section becomes:

$$\sigma_{lA \toHX} = f_A^H \otimes 0\sigma_{lH \to lH} + f_A^0 \otimes (1 f_g^H \otimes 0\hat{\sigma}_{lH \to lH} - 1 f_g^0 \otimes 0\sigma_{lH \to lH})$$  

$$= (f_A^H - f_A^0 \otimes 1 f_g^H) \otimes 0\sigma_{lH \to lH} + f_A^0 \otimes 1\hat{\sigma}_{l g \to lH}$$ (6)

4) Needless to say, colinear singularities associated with light partons are $\overline{\text{MS}}$ subtracted.

5) As mentioned earlier, we do not consider “small-x” corrections in this paper.

6) Here “unsubtracted” refers to heavy quark mass effects. As mentioned before, colinear singularities due to light partons are always subtracted as in the $\overline{\text{MS}}$ scheme.
The three terms on the right-hand-side of the first line are: heavy-flavor excitation (HE), heavy-flavor creation (HC), and the subtraction term; the last physically represents the overlap between the two production mechanisms. Both \( \sigma_{lH \rightarrow lH} \) and \( \sigma_{l \rightarrow lH \bar{H}} \) (unsubtracted) contain the full \( m_H \) dependence. For \( Q \gg m_H \), it is useful to view the right-hand side of Eq.6 as on the first line. The quantity in the parenthesis (really just \( \hat{\sigma}_{l \rightarrow lH \bar{H}} \)) is free of large \( \ln \frac{Q^2}{m_H^2} \) logarithms because of the subtraction; it is infra-red safe. The whole term remains at numerical order \( \alpha_s \times \mathcal{O}(1) \) (in contrast to the unsubtracted \( \sigma_{l \rightarrow lH \bar{H}} \) which is the LO FFN scheme result with a large log factor, \( \alpha_s \times \mathcal{O}(\ln \frac{Q^2}{m_H^2}) \)) in the large \( Q \) limit. Thus, one recovers the LO ZMVFS formula with the HE contribution as the dominant term for the cross-section. In the threshold region, \( Q \sim m_H \), it is more useful to focus attention on the second line in Eq.6: the two terms inside the parenthesis are individually small in this region and, in addition, they cancel each other up to order \( \alpha_s^2 \) since both satisfy the evolution equation (to this order) with the same boundary conditions (assuming there is no non-perturbative charm). As a consequence, the full cross-section is dominated by the second (HC) term – which is the FFN scheme result to order \( \alpha_s \).

The explicit expression for the subtraction term clearly shows how it overlaps the HE and HC mechanisms, hence leads to the appropriate cancellations in the respective kinematic regions. The same principle applies in other lepton-hadron and hadron-hadron processes to all orders of perturbation theory [17]. Numerical calculations based on Eq.6 confirm the features described above for relevant physical cross-sections. See Ref. [15] and the talk by Schmidt [18] for plots of \( F_2^c(x, Q) \) with individual contributions from the three terms on the right-hand side of Eq.6 which explicitly illustrate these features.

5 **Complementarity between HE and High-order HC Contributions**

The GM-VFN (ACOT) scheme highlights some overlapping features of HE and higher order HC mechanisms for heavy quark production (hence the need for the subtraction to avoid double counting): the two are not mutually exclusive, as sometimes perceived; rather, they are complementary. For the total inclusive cross-section (e.g. structure functions \( F_2^c(x, Q) \) in DIS), in particular, the HE contribution represents the result of resumming the collinear parts of HC diagrams to all orders in the running coupling. This provides an efficient method of obtaining important quantitative results without having to calculate many complicated higher order diagrams in the FFN scheme.

The trade-off is that some information on the differential distributions (e.g. transverse momentum spectrum of the heavy particle) is integrated over in the resummation, hence becomes less accessible. For example, if one is interested in the \( p_t \) distribution of the charm quark in DIS lepto-production with respect to the virtual photon-target axis, the HE contribution, as well as the subtraction term, in Eq.6 will be formally proportional to the delta function \( \delta(p_t) \). This is, mathematically, a distribution (rather than an ordinary function) which needs to be folded with a suitable smearing function – some combination of experimental resolution function and theoretical \( p_t \) distribution (see below) – in order to produce meaningful physical cross-sections. Away from the small \( p_t \) region, the low-order FFN scheme
diagrams will be the logical place to start in obtaining the leading contributions to the physical $p_t$ spectrum. Cf. the talk by Schmidt [18]. To obtain a precise theory of $p_t$ distributions over the entire range, a different kind of resummation will be required. Thus, for heavy quark production in PQCD, as in many quantum mechanical problems in general (say, the double-slit problem), the appropriate way to formulate the theory depends intimately on the physics question asked.

6 Recent Developments

The ACOT scheme has been applied to the analysis of recent charm production data in neutrino scattering [19], yielding the most up-to-date information on the strange quark content of the nucleon.

Recently, in the wake of new precision measurements of the total $F_2(x, Q)$ at small $x$ [20] (where charm final states consist of up to 25% of the cross-section) as well as first measurements of $F_2^c(x, Q)$ [21], it has become obvious that a more precise theoretical treatment of charm production in NC DIS than those used so far (the ZM-VFN and FFN schemes described earlier) is now needed. A first step is taken by the CTEQ group, performing a new global QCD analysis of parton distributions [22] based on the GM-VFN scheme of ACOT (to order $\alpha_s$ only). A similar study has been done by MRRS [23], using a related procedure (see also [24]). An order $\alpha_s^2$ calculation in the general mass scheme is within reach [25] since much of the known FFN scheme results to this order can be used as the starting point.

Of equal importance is an improved theoretical treatment of various heavy quark cross-sections in hadron-hadron scattering beyond the conventional calculational schemes, particularly the NLO FFN scheme results which suffer from the problems mentioned in Sec. 2. Some preliminary results have been obtained in the ACOT scheme [26,27]. The qualitative features of these results, compared to existing calculations, are similar to those on lepto-production described briefly in Sec. 4. As expected, the inclusion of heavy-flavor-excitation contributions leads to: (i) reduced dependence on the (unphysical) scale parameter $\mu$; and (ii) an increase in the predicted cross-section over most of the range of $p_t$ of the produced heavy quark, as preferred by data [11]. Detailed phenomenology of the inclusive cross-section, as well as correlations of final-state heavy particles still need to be pursued [27].

7 Concluding Remarks

The general-mass variable-flavor-number scheme of Ref. [6,14,15,17] generalizes the familiar zero-mass variable-flavor-number QCD parton framework (valid only at large $Q$ scales) to include quark mass effects, and it reproduces the results of the widely used FFN scheme heavy quark calculations (valid in the one-large-scale region $Q \sim m_H$). It naturally unifies the two contrasting conventional approaches in a well-defined renormalization and factorization scheme. First results on lepto- and hadro-production demonstrate improvements over existing calculations both in smaller $\mu$-dependence and in increased cross-sections.
However, the more complete theory is not a cure-for-all. Although detailed phenomenology has yet to be done, there is no doubt that further developments in several directions are needed to reach a full understanding of heavy quark production in QCD. For instance, (i) the general mass scheme should be implemented to higher orders for both lepto- and hadro-production; (ii) the significance of possible large logarithms of the type $\ln(Q^2/s), \ln(m_H^2/s) \sim \ln x$ – the small-$x$ problem – needs to be better understood [12]; (iii) both perturbative and non-perturbative aspects of the hadronization of the heavy quarks deserve further study; and (iv) the question “Are there non-perturbative charm/bottom components inside hadrons?” needs to be answered. The last question, having been discussed in the literature for many years, has to be carefully investigated phenomenologically; and this can be done only in the framework of the general-mass scheme, since the existence of non-perturbative non-zero-mass parton is excluded by assumption in the FFN scheme.

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