PQCD ANALYSIS OF INCLUSIVE HEAVY HADRONS DECAYS

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We develop the perturbative QCD formalism for inclusive heavy hadron decays. Transverse degrees of freedom of partons are introduced to facilitate the factorization of the heavy hadron decays.

In this talk, we shall derive the PQCD factorization formula for the semileptonic decay $B \to X_u \ell \nu$ in a rigorous way, to demonstrate how to use techniques of PQCD for heavy hadron decays. The factorization procedures demand the inclusion of the transverse degrees of freedom of partons at the end point of the charge lepton spectrum where an outgoing jet is present. Hence, we have to perform the resummation of large perturbative corrections in the transversal configuration space using the technique developed in [1], which is also accurate up to next-to-leading logarithms.

We work in the rest frame of the $B$ meson, and choose the following light-cone components for relevant momenta,

\[ P_B = (P_B^+, P_B^-, 0, 0), \quad p_\ell = (p_\ell^+, 0, 0, 0), \quad p_\nu = (p_\nu^+, p_\nu^-, p_\nu^\perp), \]

with $P_B^+ = P_B^- = M_B/\sqrt{2}$ and $p_\nu^2 = 0$. The independent variables are identified as $p_\ell^+$, $p_\ell^-$ and $p_\nu^+$, and their relations to $E_\ell$, $q^2$ and $q_0$ are $E_\ell = p_\ell^+/\sqrt{2}$, $q^2 = 2p_\ell^+p_\nu^-$, and $q_0 = (p_\ell^+ + p_\nu^+ + p_\nu^-)/\sqrt{2}$, respectively. We define $P_b = P_B - p$ as the $b$ quark momentum, which satisfies $P_b^2 \approx M_b^2$, $M_b$ being the $b$ quark mass. $p$ is the kicks from the light components inside the $B$ meson, which has a large plus component $p_\ell^+$ and small transverse components $p_\nu^\perp$. The $b$ quark decays into a $u$ quark with momentum $P_u = P_B - p - q$. We have distinguished the $B$ meson momentum $P_B$ from the $b$ quark momentum $P_b$ here.

It is more convenient to employ the scaling variables

\[ x = \frac{2E_\ell}{M_B}, \quad y = \frac{q^2}{M_B^2}, \quad y_0 = \frac{2q_0}{M_B}, \]

instead of the dimensionful ones $E_\ell$, $q^2$ and $q_0$. Note that the scaling variables are defined in terms of the $B$ meson mass $M_B$, since we formulate the factorization according to the $B$ meson kinematics. For massless leptons, it is easy to show, using the momentum configurations defined in eq. (1), that the phase
space is given by

\[ 0 \leq x \leq 1, \quad 0 \leq y \leq x, \quad \frac{y}{x} + x \leq y_0 \leq y + 1. \tag{3} \]

In the end point region with \( x \to 1 \) \((p^+_\ell \to M_B/\sqrt{2})\) and \( y \to 0 \) \((p^-_\nu \to 0)\), we have \( y_0 \to 1 \) \((p^+_u \to 0)\) and \( p \to 0 \). The \( u \) quark then has a large minus component \( P^-_u = (1 - y/x)M_B/\sqrt{2} \) but a very small plus component \( P^+_u = (1 - y_0 - y/x)M_B/\sqrt{2} \), and thus a very small invariant \( P^2_u = M^2_B(1 - y_0 + y) \), which forms an on-shell jet subprocess. The \( u \) quark travels a long distance of \( \mathcal{O}(1/\Lambda_{QCD}) \) before hadronized. Besides, the \( B \) meson is dominated by soft dynamics, which is the origin of the soft function stated in the Introduction. The remaining dominant subprocess is the hard one, which contains the weak decay vertices. Therefore, the important contributions are factorized into the soft \( (S) \), jet \( (J) \) and hard \( (H) \) subprocesses.

The factorization formula for the inclusive semileptonic decay \( B \to X_u \ell\nu \) is written as

\[
\frac{1}{\Gamma^{(0)}_\ell} \frac{d^3\Gamma}{dxdydy_0} = M^3_B \int_{z_{\text{min}}}^{z_{\text{max}}} dz \int d^2p_\perp \times S(z, p_\perp, \mu)J(z, P^-_u, p_\perp, \mu)H(z, P^-_u, p_\perp, \mu), \tag{4} \]

with the momentum fraction \( z \) defined by \( z = P^+_B / P^+_B = 1 - p^+ / P^+_B \) and \( \Gamma^{(0)}_\ell = \frac{G^2_F}{16\pi}|V_{ub}|^2 M^2_B \). \( \mu \) in eq. (3) is the renormalization and factorization scale. The triple differential decay rate is, of course, \( \mu \) independent. Note that in the region \( y \to x \sim 1 \) the outgoing \( u \) quark becomes soft and eq. (4) fails. We shall show that contributions from this dangerous region are suppressed by phase space. The upper limit of \( z \) takes the value \( z_{\text{max}} = 1 \) in our analysis. If performing the factorization according to the \( b \) quark kinematics, one must assume \( z_{\text{max}} = M_B/M_b \), which is greater than 1, in order to fill up the kinematic window. It has been explained that \( z_{\text{max}} > 1 \) is not allowed in perturbation theory, and is thus of nonperturbative origin. From the kinematic constraints in eq. (3) and the on-shell condition of the \( u \)-quark jet, the lower limit of \( z \) should be \( z_{\text{min}} = x \), instead of \( z_{\text{min}} = 0 \).

The tree-level expressions for the convolution factors \( J \) and \( H \) are given by

\[
J^{(0)} = \delta(P^2_u) = \delta \left( M^2_B \left[ 1 - y_0 + y - (1 - z)(1 - y/x) - \frac{2p_\perp \cdot p_\perp}{M^2_B} - \frac{p_\perp^2}{M^2_B} \right] \right),
\]

\[
H^{(0)} \propto (P_b \cdot p_\nu)(p_\ell \cdot P_u) = ((P_B - p) \cdot p_\nu)(p_\ell \cdot P_u)
\]

\[
\propto (x - y) \left( y_0 - x - (1 - z)\frac{y}{x} + \frac{2p_\perp \cdot p_\perp}{M^2_B} \right). \tag{5} \]
Equation (4) can be regarded as an expression at the intermediate stage in the derivation of conventional factorization theorems. If the $p_{\perp}$ dependence in $J$ and $H$ is negligible, the variable $p_{\perp}$ in $S$ can be integrated over, and eq. (4) reduces to the conventional factorization formula. However, it is obvious from eq. (5) that at least the $p_{\perp}$ dependence in $J$ is not negligible, especially in the end-point region. This is the reason we introduce the transverse degrees of freedom into our analysis.

Suppose we consider higher-order corrections to eq. (4) from a gluon crossing the final state cut, and route the loop momentum $\ell$ through, say, the jet subprocess. Without losing generality, we approximate $J(p^+ + \ell^+, P_u^-, P_{u+} + \ell_{\perp}) \approx J(p^+, P_u^-, P_{u+} + \ell_{\perp})$ according to the kinematic relations $\ell^+ < p^+$, $\ell^- < P_u^-$, and $\ell_{\perp} \approx p_{\perp}$. Hence, the loop integral cannot be performed unless the dependence of $J$ on transverse momentum is known. This difficulty can be removed by Fourier transform,

$$J(p^+, P_u^-, P_{\perp} + \ell_{\perp}) = \int \frac{d^2 b}{(2\pi)^2} \tilde{J}(p^+, P_u^-, b) e^{i(\ell_{\perp} \cdot b)}$$

where the impact parameter $b$ (Fourier conjugate variable of $p_{\perp}$) measures the transverse distance travelled by the jet. Using eq. (6), the $\ell_{\perp}$ dependence is decoupled from the jet function, and the factor $e^{i(\ell_{\perp} \cdot b)}$ is absorbed into the loop integral, which can then be performed. Therefore, an extra factor $e^{i(\ell_{\perp} \cdot b)}$ is associated with each gluon crossing the final state cut in our formalism.

To further simplify the factorization formula, we neglect those terms involving $p_{\nu_{\perp}}$ in $J$ and $H$. This is a good approximation for $x \to 0$ and $x \to 1$, since for $x \to 0$ contributions from transverse momenta are not important, and for $x \to 1$ the magnitude $p_{\nu_{\perp}} = M_B \sqrt{y_0 - x - y/x}$ vanishes. Equation (4) then becomes

$$\frac{1}{\Gamma^{(0)}_{\ell}} \frac{d^3 \Gamma}{dx dy y_0} = M_B^2 \int_x^1 dz \int \frac{d^2 b}{(2\pi)^2} \tilde{S}(z, b, \mu) \tilde{J}(z, P_u^-, b, \mu) \tilde{H}(z, P_u^-, \mu).$$

It can be shown that the dominant subprocesses $J$, $S$ and $H$ contain large logarithms from radiative corrections. In particular, $J$ gives rise to double (leading) logarithms in the end point region. These large corrections spoil the perturbation theory and must be organized. One has to resum these large corrections up to next-to-leading logarithms. The first step in resummation is to map out the leading regions of radiative corrections. In the collinear region with the loop momentum $\ell$ parallel to $P_u$ and in the soft region with $\ell \to 0$ we
can eikonalize the heavy b-quark line. With the eikonal approximation, the b-quark propagator is expressed as $1/v \cdot \ell$ to the order $1/M_B$ with $v = (1, 1, 0_\perp)$. Hence, the factorization in eq. (5) is in fact valid to $O(1/M_B)$. The physics involved in this approximation is that a soft gluon or a gluon moving parallel to $P_u$ can not explore the details of the b quark, and its dynamics can be factorized. This is consistent with the HQEFT, where the b quark is treated as a classical relativistic particle carrying color source. Since the large mass $M_B$ does not appear in the eikonal propagator, the only large scale in $J$ is $P_u^\perp$. After eikonalization procedures, one obtain the Sudakav exponent that resumes large double logarithms arise from both collinear and soft regions,

$$S = 2 \int_{1/b}^{m_\mu} \frac{dp}{p} \int_{1/b}^{\mu} \frac{d\mu}{\mu} A(\alpha_s(\mu)) ,$$

$$A = C_F \frac{\alpha_s}{\pi} + \left[ \frac{67}{9} - \frac{\pi^2}{3} - \frac{10}{27} n_f \right] (\frac{\alpha_s}{\pi})^2 ,$$

with $C_F = 4/3$ the color factor.

Having summed up the double logarithms, we concentrate on the single logarithms in $\tilde{S}, H$ and the initial condition $\tilde{J}(b, \mu)$. Since both the differential decay rate and the Sudakov exponent $s(P_u^-, b)$ are RG invariant, we have the following RG equations:

$$D \tilde{J}(b, \mu) = -2\gamma_q \tilde{J}(b, \mu) ,$$

$$D \tilde{S}(b, \mu) = -\gamma_S \tilde{S}(b, \mu) ,$$

$$D H(P_u^-, \mu) = (2\gamma_q + \gamma_S) H(P_u^-, \mu) ,$$

with

$$D = \mu \frac{\partial}{\partial \mu} + \beta(g) \frac{\partial}{\partial g} .$$

$\gamma_q = -\alpha_s/\pi$ is the quark anomalous dimension in axial gauge, and $\gamma_S = -(\alpha_s/\pi)C_F$ is the anomalous dimension of $\tilde{S}$. Integrating eq. (10), we obtain the evolution of all the convolution factors,

$$\tilde{J}(z, P_u^-, b, \mu) = \exp \left[ -2s(P_u^-, b) - 2 \int_{1/b}^{\mu} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_q(\alpha_s(\bar{\mu})) \right] \tilde{J}(z, b, 1/b) ,$$

$$\tilde{S}(z, b, \mu) = \exp \left[ - \int_{1/b}^{\mu} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_S(\alpha_s(\bar{\mu})) \right] f(z, b, 1/b) ,$$

$$H(z, P_u^-, \mu) = \exp \left[ - \int_{\mu}^{P_u^-} \frac{d\bar{\mu}}{\bar{\mu}} [2\gamma_q(\alpha_s(\bar{\mu})) + \gamma_S(\alpha_s(\bar{\mu}))] \right] H(z, P_u^-, P_u^+) .$$
We shall neglect the intrinsic $b$ dependence of the distribution function $f$ below. If Sudakov suppression in the large-$b$ region is strong, we may drop the evolution of $f$ and $\tilde{J}$ in $b$, which is proportional to $\alpha_s(1/b)$. Hence, we assume $f(z, b, 1/b) = f(z), \tilde{J}(z, b, 1/b) = \tilde{J}^{(0)}(z, b)$, the Fourier transform of the tree-level expression in eq. (5), and $H(z, P_u^-, P_u^-) = H^{(0)}(z, P_u^-)$.

Substituting eq. (12) into (7), we derive the factorization formula for the inclusive semileptonic $B$ meson decay,

$$\frac{1}{\Gamma_t^{(0)}} \frac{d^3\Gamma}{dxdyd\theta} = M_B^2 \int_x^1 dz \int_0^\infty db \frac{b f(z)\tilde{J}^{(0)}(z, b)H^{(0)}(z, P_u^-)}{2\pi} \times \exp[-S(P_u^-, b)].$$

The complete Sudakov exponent $S$ is given by

$$S(P_u^-, b) = 2s(P_u^-, b) - \frac{5}{3\beta_1} \ln\frac{\hat{P}_u^-}{b},$$

with $\hat{P}_u^- = \ln(P_u^-/\Lambda)$, which combines all the exponents in eq. (12) and includes both leading and next-to-leading logarithms. It is straightforward to observe from eq. (14) that the Sudakov form factor $e^{-S}$ falls off quickly at large $b \sim 1/\Lambda$, where $\alpha_s(1/b) > 1$ and perturbation theory fails. Hence, the Sudakov form factor guarantees that main contributions to the factorization formula come from the small $b$, or short-distance, region, and the perturbative treatment is indeed self-consistent. We stress that our formalism is applicable to the entire range of the spectrum.

Since all the double logarithmic corrections have been absorbed into the jet subprocess, the soft function $\tilde{S}$ contains only soft single logarithms. These single logarithms can be summed by solving the RG equation $\mathcal{D}\tilde{S} = -\gamma S \tilde{S}$ as shown in eq. (10). The solution has been given in eq. (12),

$$\tilde{S}(z, b, \mu) = \exp \left[ - \int_{1/b}^\mu \frac{d\mu}{\mu} \gamma_S(\alpha_s(\mu)) \right] f(z, b),$$

where the initial condition $f(z)$ for the evolution of $\tilde{S}$ must be determined phenomenologically. From the definition of $\tilde{S}$, it is obvious that $f$ depends only on the properties of the bound state $|B\rangle$, but not on the particular short distance subprocess. Therefore, $f$ is a process-independent universal function describing the distribution of the $b$ quark inside a $B$ meson. Since the intrinsic $b$ dependence (the perturbative $b$ dependence has been collected into the exponent $S$) is not known yet, we take the ansatz $f(z, b) = f(z)\exp(-\Sigma(b))$. 

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which leads to $\Sigma \to 0$ as $b \to 0$ according to the definition $f(z, b = 0) \equiv f(z)$. It is also natural to assume $\Sigma > 0$ for all $b$ from the viewpoint that the $b$ quark is bounded inside the $B$ meson. Hence, the intrinsic $b$ dependence provides further suppression. The nonperturbative function $f(z)$, identified as the $B$-meson distribution function, can be expressed as the matrix element of the $b$ quark fields, whose first three moments are:

\[
\begin{align*}
\int_0^1 f(z)dz &= 1, \\
\int_0^1 f(z)(1-z)dz &= \frac{\bar{\Lambda}}{m_B} + O(\Lambda_{\text{QCD}}^2/m_B^2), \\
\int_0^1 f(z)(1-z)^2dz &= \frac{1}{m_B^2} \left(\lambda^2 - \frac{\lambda_1}{3}\right) + O(\Lambda_{\text{QCD}}^3/m_B^3). \quad (16)
\end{align*}
\]

Though the exponent $\Sigma$ is unknown, we can, however, extract its leading behavior by means of the IR renormalon analysis. Note that the perturbative Sudakov factor $e^{-\Sigma}$ in Eq. (15) becomes unreliable as the transverse distance $b$ approaches $1/\Lambda_{\text{QCD}}$. Near this end point, $\alpha_s(1/b)$ diverges, and IR renormalon contributions are significant. We reexpress the RG result of the evolution of the distribution function, which is contained in the second term of $S$, as:

\[
W = \exp \left[4\pi C_F \int \frac{d^4l}{(2\pi)^4} \frac{v_\mu v_\nu}{(v \cdot l)^2} 2\pi \delta(l^2) \alpha_s(l_T^2) e^{i l_T \cdot b} N^{\mu\nu} \right]. \quad (17)
\]

The loop integral corresponds to the correction from a real soft gluon attaching the two valence $b$ quarks, whose propagators have been replaced by the eikonal lines in the direction $v = (1, 1, 0)$. The tensor $N^{\mu\nu} = g^{\mu\nu} - (n^\mu l^\nu + l^\mu n^\nu)/(n \cdot l) + n^2 l^{\mu\nu}/(n \cdot l)^2$ comes from the gluon propagator in axial gauge $n \cdot A = 0$. We have set the argument of the running $\alpha_s$ to $l_T^2$, which is conjugate to the scale $b$ of the distribution function.

Substituting the identity $\alpha_s(l_T^2) = \pi \int_0^\infty d\sigma \exp[-2\sigma \beta_1 \ln(l_T/\Lambda_{\text{QCD}})], \beta_1 = (33 - 2n_f)/12$, into Eq. (17), and performing the loop integral for $n \propto (-1, 1, 0)$, we obtain:

\[
W = \exp \left[C_F \int_0^\infty d\sigma \left(\frac{b\Lambda_{\text{QCD}}}{2}\right)^{2\sigma \beta_1} \frac{\Gamma(-\sigma \beta_1)}{\Gamma(1 + \sigma \beta_1)} \right]. \quad (18)
\]

It is easy to observe that the pole of $\Gamma(-\sigma \beta_1)$ at $\sigma \to 0$ gives the perturbative anomalous dimension of the distribution function appearing in Eq. (8). The extra poles at $\sigma \to 1/\beta_1, 2/\beta_1, \ldots$ then correspond to the IR renormalons, giving
corrections of powers $b^2$, $b^4$, ..., respectively. These renormalons generate singularities, which must be compensated by the nonperturbative power corrections in order to have a well-defined perturbative expansion. Using the "minimal" ansatz of picking up only the leading (left-most) renormalon contribution, we parametrize the exponent $\Sigma(b)$ by $\Sigma(b) = c'm_b^2 b^2$, corresponding to the fact that the power corrections start at $O(b^2)$. Certainly, other parametrizations consistent with this fact are equally good. We postulate the following two-parameter distribution function for the $B$ meson,

$$f_B(z) = N \frac{z(1-z)^2}{((z-a)^2 + \epsilon z)^2} \theta(1-z), \quad (19)$$

The charged lepton spectrum for the decay $B \to X_u \ell \nu$ from the naive quark model is obtained by simply choosing $f(z) = \delta(1-z)$ and ignoring the transverse momentum dependence in $J(0)$ and $H(0)$. A simple calculation leads to

$$\frac{1}{\Gamma_\ell^{(0)}} \frac{d\Gamma}{dx} = \frac{x^2}{6} (3 - 2x). \quad (20)$$

This spectrum does not fall off at the end point of the spectrum, contradicting the observed behavior of the inclusive semileptonic decays of $B$ mesons. The discrepancy implies that the tree-level analysis is not appropriate, especially in the end-point region where PQCD corrections are important.

The spectrum from the parton model without Sudakov suppression is obtained by adopting $H^{(0)} = (x-y)[y_0 - x - (1-z)y/x]$ and $P_u^2 = M_B^2 [1 - y_0 + y - (1-z)(1-y/x)]$. With integration over $y_0$, we derive

$$\frac{1}{\Gamma_\ell^{(0)}} \frac{d\Gamma}{dx} = \int_0^x dy \int_x^1 dz f(z)(x-y)(y+z-x). \quad (21)$$

At last, including Sudakov suppression and all other single logarithms, we arrive at the charged lepton spectrum of the $B \to X_u \ell \nu$ decay that takes into account both large perturbative and nonperturbative corrections,

$$\frac{1}{\Gamma_\ell^{(0)}} \frac{d\Gamma}{dx} = M_B \int_0^x dy \int_0^{1/A} db \int_x^1 dz f(z)(x-y)\xi \left[ (z+y-x)J_1(\xi M_B b) - \frac{2}{M_B^2} \xi J_2(\xi M_B b) + \xi^2 J_3(\xi M_B b) \right] e^{-S(P_u^{-},b)}, \quad (22)$$

with $\xi = \sqrt{(x-y)(z/x - 1)}$ and $J_1, J_2, J_3$ are the Bessel functions of order 1, 2 and 3 respectively.
We stress that our results do not violate the conclusion from HQEFT, if they were interpreted in a proper way. To confirm this, we identify 
\[ \mathbf{P}_b = \left( \frac{M_b^2}{2M_B^2} P_b, P_B, 0 \right) \]
as the momentum carried by a free \( b \) quark in factorization theorems for \( B \) meson decays, where the minus component \( P_{b-} \) has been set to \( P_B^- \). That is, the free \( b \) quark is not at rest inside the \( B \) meson. We then reexpress eq. (21) into a form similar to that in (3):

\[
\frac{1}{\Gamma_{\ell}^{(0)}} \frac{d\Gamma}{dx} = F(x) \theta \left( \frac{M_b^2}{M_B^2} - x \right) + F \left( \frac{M_b^2}{M_B^2} \right) M(x) , \tag{23}
\]

with \( F(x) = x^2(3 - 2x)/6 \) being the quark-model prediction derived from the conventional approaches and

\[
F \left( \frac{M_b^2}{M_B^2} \right) M(x) = \int_0^x dy \int_x^1 dz f(z)(x-y)(y+z-x) - F(x) \theta \left( \frac{M_b^2}{M_B^2} - x \right) . \tag{24}
\]

The step function in eq. (23) specifies the maximal \( E_\ell \) in the decay of a free \( b \) quark with the above momentum \( \mathbf{P}_b \). The function \( M(x) \), representing nonperturbative corrections to the \( b \) quark decay, coincides with the shape function \( S(x) \) defined in (3).

We shall show that the contribution from \( M(x) \) to the total decay rate is indeed of \( \mathcal{O}(1/M_B^2) \). Integrating eq. (24) over \( x \), we obtain

\[
F \left( \frac{M_b^2}{M_B^2} \right) \int_0^1 M(x) dx = \frac{1}{12} \int_0^1 dz \int_x^1 f(z)(x-y)(y+z-x) - \int_0^1 dz \left( \frac{M_b^2}{M_B^2} - x \right) . \tag{25}
\]

An arbitrary structure function \( f \), which possesses the same moment as in eq. (16), can be expanded in terms of \( \delta \)-functions:

\[
f(z) = \delta(1-z) - \frac{\bar{\Lambda}}{M_B} \delta'(1-z) + \mathcal{O}(\bar{\Lambda}^2/M_B^2) . \tag{26}
\]

Inserting eq. (26) into (24), we justify straightforwardly that the nonperturbative correction

\[
F \left( \frac{M_b}{M_B} \right) \int_0^1 M(x) dx = \int_{M_b^2/M_B^2}^1 F(x) dx - \frac{1}{12} \frac{\bar{\Lambda}}{M_B} \int_0^1 dz \int_x^1 f(z)(x-y)(y+z-x) + \mathcal{O}(\bar{\Lambda}^2/M_B^2)
\]

\[
= \frac{1}{3} \frac{\bar{\Lambda}}{M_B} - \frac{1}{3} \frac{\bar{\Lambda}}{M_B} + \mathcal{O}(\bar{\Lambda}^2/M_B^2) \tag{27}
\]

vanishes at \( \mathcal{O}(1/M_B) \) as concluded in (3). In summary, the quark-model contribution from the window between \( x = M_b^2/M_B^2 \) and \( x = 1 \), with a width of
\(\mathcal{O}(1/M_B)\), cancels the \(\mathcal{O}(1/M_B)\) correction from the structure function, such that the nonperturbative correction is of \(\mathcal{O}(1/M_B^2)\). Our result provides an explicit dynamical demonstration the validity of global quark-hadron duality.

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