Subleading Shape-Function Effects and the Extraction of $|V_{ub}|$

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Abstract

We derive a class of formulae relating moments of $B \to X_u \ell \bar{\nu}$ to $B \to X_s \gamma$ in the shape function region, where $m_X^2 \sim m_b \Lambda_{QCD}$. We also derive an analogous class of formulae involving the decay $B \to X_s \ell^+ \ell^-$. These results incorporate $\Lambda_{QCD}/m_b$ power corrections, but are independent of leading and subleading hadronic shape functions. Consequently, they enable one to determine $|V_{ub}|/|V_{tb}V_{ts}^*|$ to subleading order in a model-independent way.

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I. INTRODUCTION

The study of decays of the $B$ meson allows us to probe QCD and flavour physics. The program's goals include, on the one hand, precision measurements of Standard Model parameters and, on the other hand, searches for new physics. Short-distance physics is encoded in Wilson coefficients of local operators. By comparing measurements of these coefficients with theoretical predictions, signals of new physics may be found. High sensitivity to new physics is provided by the so-called rare decays, namely those channels involving flavour-changing neutral currents, since they do not occur at tree level in the Standard Model. Measurements of the inclusive rare process $B \rightarrow X_s \gamma [1, 2, 3, 4, 5]$ have provided significant constraints on extensions to the Standard Model. The more complicated decay $B \rightarrow X_s \ell^+ \ell^-$ is complementary to $B \rightarrow X_s \gamma$, as its effective Hamiltonian includes two extra operators. Moreover, additional observables are available, such as the $q^2$ spectrum and the forward-backward asymmetry, which have been the focus of much work. Recently, it was noted that an angular decomposition provides a third observable sensitive to a different combination of Wilson coefficients [6]. Belle and BABAR have already made initial measurements of $B \rightarrow X_s \ell^+ \ell^-$ [7, 8].

Precision measurements also provide determinations of elements of the CKM matrix or, equivalently, the angles and sides of the unitarity triangle. By overconstraining these, the flavour structure of the Standard Model is subjected to rigorous examination. For the decay $B \rightarrow X_c \ell \bar{\nu}$, experimental and theoretical uncertainties are under control, and consequently $|V_{ub}|$ is one of the best-determined elements of the CKM matrix. From $B \rightarrow X_u \ell \bar{\nu}$, we can also determine $|V_{ub}|$ [9, 10, 11, 12].

However, inclusive $B$ decays often require a trade-off between theoretical and experimental difficulty: if phase-space cuts are necessary experimentally, then the spectra will be less inclusive and the corresponding theory more complicated. In this respect, $B \rightarrow X_c \ell \bar{\nu}$ and $B \rightarrow X_u \ell \bar{\nu}$ are markedly different. The former is sufficiently inclusive to enable the use of a local operator product expansion (OPE) [13], in which non-perturbative corrections appear as an expansion in inverse powers of $m_b$. This formalism has been calculated to order $1/m_b^3$ [14] (and recently to order $1/m_b^4$ [15]), with the relevant non-perturbative matrix elements defined via the Heavy Quark Effective Theory (HQET) [16, 17, 18]. In contrast, in $B \rightarrow X_u \ell \bar{\nu}$ experimental cuts (e.g. cuts on $E_\ell$ or $m_{X_u}^2$) are required in order to eliminate the dominant $b \rightarrow c$ background. In many cases, we are restricted to a region in which $m_{X_u}^2 \sim m_b \Lambda_{QCD}$ and the local OPE breaks down. In this so-called endpoint or shape function region [19], the set of outgoing hadronic states becomes jet-like and the relevant degrees of freedom are collinear and ultrasoft modes. The Soft-Collinear Effective Theory (SCET) [20, 21, 22, 23] is then a powerful theoretical method.

Similarly, $B \rightarrow X_s \gamma$ measurements employ a cut on the photon energy. In Refs. [24, 25] it was shown that the shape function region is also relevant for $B \rightarrow X_s \ell^+ \ell^-$. Here, cuts are made in the dileptonic mass spectrum to remove the largest $c\bar{c}$ resonances, namely the $J/\Psi$ and $\Psi'$. These leave two perturbative windows, the low-$q^2$ and high-$q^2$ regions. At low $q^2$, where the rate is higher, an additional cut is needed: a hadronic invariant-mass cut is imposed in order to eliminate the background $b \rightarrow c(\rightarrow s \ell^+ \nu) \ell^- \bar{\nu}$.

At leading order (LO) in $\Lambda_{QCD}/m_b$, decay rates now depend upon a non-perturbative,
and hence analytically incalculable, shape function. However, this function is process-independent and appears in both \( B \to X_u \ell \bar{\nu} \) and \( B \to X_s \gamma \), for example. One can thus measure the leading-order shape function from the photon energy spectrum of \( B \to X_s \gamma \) and use the result in the \( B \to X_u \ell \bar{\nu} \) spectrum, or, more directly, express the semileptonic rate in terms of the radiative rate instead of the shape function [26, 27, 28, 29]. In this way, model dependence can be avoided in the determination of \(|V_{ub}|\).

At subleading order, the situation is far more complicated, with several universal shape functions occurring in different combinations [30, 31, 32, 33, 34, 35]. In this paper, we construct combinations of shape-function-dependent decay rates that are protected from non-perturbative effects to second order in the power expansion. Through this procedure, we obtain formulae for \(|V_{ub}|/|V_{ub}V_{ts}^*|\) that are free from the hadronic uncertainties arising from the leading and subleading shape functions. This method uses moments of the fully differential decay spectra of \( B \to X_u \ell \bar{\nu} \) and \( B \to X_s \gamma \) (and, optionally, \( B \to X_s \ell^+ \ell^- \)).

The rest of this paper is organized as follows. In Sec. 2, together with Appendices A and B, we present the basic formalism needed for our work. This includes power corrections for the triply differential decay spectra of the semileptonic processes and the photon energy spectrum of \( B \to X_s \gamma \). In Sec. 3, we derive and discuss our results, eliminating shape functions from expressions for \(|V_{ub}|\) at next-to-leading order (NLO). We conclude in Sec. 4.

II. FORMALISM

In this section, we briefly review the formalism and results from Refs. [24, 32, 36] that we shall use in this paper (see these references for further details).

The inclusive decay rate for \( \bar{B} \to X_u \ell \bar{\nu} \) (\( \bar{B} \to X_s \gamma \)) is proportional to \( W_{\mu \nu} L_{\mu \nu} \), where \( L_{\mu \nu} \) is the leptonic (photonic) tensor and \( W_{\mu \nu} \) is the hadronic tensor, which can be written as

\[
W_{\mu \nu} = \frac{1}{2m_B} \sum_X (2\pi)^3 \delta^4(p_B - q - p_X) \langle \bar{B} | J_{\mu}^\dagger | X \rangle \langle X | J_{\nu} | \bar{B} \rangle = -g_{\mu \nu} W_1 + v^\mu v^\nu W_2 + i \epsilon_{\mu \rho \alpha \beta} v^\alpha q^\beta W_3 + q_\mu q_\nu W_4 + (v_\mu q_\nu + v_\nu q_\mu) W_5.
\]

Here, \( v^\mu \) is the velocity of the \( B \) meson and \( q^\mu \) is the \( \ell \bar{\nu} \) (\( \gamma \)) momentum. We use the hadronic current \( J_\mu = \bar{u} \gamma_\mu p_t b \) for \( B \to X_u \ell \bar{\nu} \) and relativistic normalization for the \(|\bar{B}\rangle\) states. Similarly, the inclusive decay rate for \( \bar{B} \to X_s \ell^+ \ell^- (p^+) \ell^- (p^-) \) is proportional to \( (W_{\mu L}^L L_{\mu L}^L + W_{\mu R}^R L_{\mu R}^R) \), where \( L_{\mu (L,R)} = 2 [p_+^\mu p_+^\nu + p_-^\mu p_-^\nu - g_{\mu \nu} p_+ \cdot p_- + i \epsilon_{\mu \nu \alpha \beta} p_+^\alpha p_-^\beta] \) and \( W_{\mu (L,R)} \) can be defined analogously to Eq. (1), in terms of a current \( J_{\mu (L,R)} \) [37].

Contracting \( L_{\mu \nu} \) with \( W_{\mu \nu} \) and neglecting the mass of the leptons gives the differential decay rates

\[
\frac{d\Gamma^s}{dE_{\gamma}} = \Gamma_0^s \frac{8E_\gamma}{m_B^3} (4W_1^s - W_2^s - 2E_\gamma W_5^s),
\]

\[
\frac{d^3\Gamma^u}{dE_{\ell}dq^2dE_{\nu}} = \Gamma_0^u \frac{96}{m_B^3} \left[ q^2 W_1^u + (2E_\ell E_{\nu} - q^2/2) W_2^u + q^2 (E_\ell - E_{\nu}) W_3^u \right] \theta(4E_\ell E_{\nu} - q^2),
\]

\[
\frac{d^3\Gamma^{\ell\ell}}{dq^2dE_+dE_-} = \Gamma_0^{\ell\ell} \frac{96}{m_B^3} \left[ q^2 W_1^{\ell\ell} + (2E_- E_+ - q^2/2) W_2^{\ell\ell} + q^2 (E_- - E_+) W_3^{\ell\ell} \right] \theta(4E_- E_+ - q^2),
\]
for $B \rightarrow X_s \gamma$, $B \rightarrow X_s \ell \bar{\nu}$ and $B \rightarrow X_s \ell^+ \ell^−$, respectively, where $W_{1(2)}^{\ell \ell} = W_{1(2)}^L + W_{1(2)}^R$, $W_3^{\ell \ell} = W_3^L - W_3^R$, $W_i = W_i(q^2, v\cdot q)$ and the normalization factors are

$$
\Gamma_0^s = \frac{G_F^2 m_B^3}{12 \pi^4} |V_{tb} V_{ts}^*|^2 \alpha_{em} |m_b(m_b)|^2 |C_7^{\ell}(m_b)|^2, \quad \Gamma_0^u = \frac{G_F^2 m_B^5}{192 \pi^3} |V_{ub} V_{ts}^*|^2.
$$

In SCET, it is natural to use light-cone coordinates, where we introduce vectors $n$ and $\bar{n}$ such that $n^2 = \bar{n}^2 = 0$ and $n \cdot \bar{n} = 2$. A four-vector then has components $(p^+, p^−, p_\perp) = (n \cdot p, \bar{n} \cdot p, p^\mu_\perp)$. In the region of interest, the set of hadronic states $X$ is jet-like, i.e. $p_X^+ \ll p_X^−$. For convenience we define the dimensionless variables

$$
x^\gamma_H = \frac{2E_\gamma}{m_B}, \quad x_H = \frac{2E_\ell}{m_B}, \quad \bar{y}_H = \frac{\bar{n} \cdot p_X}{m_B}, \quad u_H = \frac{n \cdot p_X}{m_B}.
$$

In terms of these variables, the decay rates are

$$
d\Gamma^s = \frac{2x_H^2}{m_B} \left\{ 4W_1^s - W_2^s - m_B x_H W_5^s \right\},
$$

$$
\frac{1}{\Gamma_0^u} \frac{d^3\Gamma^u}{dx_H d\bar{y}_H du_H} = 24m_B(\bar{y}_H - u_H) \left\{ (1-u_H)(1-\bar{y}_H)W_1^u + \frac{1}{2}(1-x_H - u_H)(x_H + \bar{y}_H - 1)W_2^u + \frac{m_B}{2} (1-u_H)(1-\bar{y}_H)(2x_H + u_H + \bar{y}_H - 2)W_3^u \right\},
$$

$$
\frac{1}{\Gamma_0^{\ell \ell}} \frac{d^3\Gamma^{\ell \ell}}{dx_H d\bar{y}_H du_H} = 24m_B(\bar{y}_H - u_H) \left\{ (1-u_H)(1-\bar{y}_H)W_1^{\ell \ell} + \frac{1}{2}(1-x_H - u_H)(x_H + \bar{y}_H - 1)W_2^{\ell \ell} + \frac{m_B}{2} (1-u_H)(1-\bar{y}_H)(2x_H + u_H + \bar{y}_H - 2)W_3^{\ell \ell} \right\},
$$

where $W_i = W_i(u_H, \bar{y}_H)$. The full phase-space limits are given in Table II of Ref. [32].

The optical theorem relates the $W_i$ to forward-scattering amplitudes, which can be calculated by taking time-ordered products of currents. An important part of the analysis is the separation of short- and long-distance contributions. The results, known as factorization theorems, may be written schematically in the form

$$
d\Gamma = H \times J \otimes f,
$$

where $\otimes$ denotes a convolution. The hard ($H$) and jet ($J$) functions encode perturbative corrections that appear at two different scales, $\mu_0 \sim m_b$ and $\mu_i \sim \sqrt{m_b \Lambda_{QCD}}$ respectively, whereas the shape function ($f$) represents non-perturbative physics.

SCET involves a power expansion in the small parameter $\lambda = \sqrt{\Lambda_{QCD}/m_b}$. At leading order in $\lambda$, rates depend on one shape function, which we denote by $f^{(0)}$.

$$
W_i^{(0)} = h_i(p_X^-, m_b, \mu) \int_0^{p_X^+} dk^+ J^{(0)}(p_X^-, k^+, \mu) f^{(0)}(k^+ + \bar{\Lambda} - p_X^+, \mu),
$$

where $\bar{\Lambda} = m_B - m_b + (\lambda_1 + 3\lambda_2)/(2m_b) + \ldots$. The first subleading shape functions occur at order $\lambda^2$ and we denote these by $f^{(2)}_{0,2}$, $f^{(4)}_{3,4}$ and $f^{(6)}_{5,6}$. These are common to the three decays,
At one-loop order, we have \( \omega \) where a prime denotes a derivative, as well as

\[
\mathcal{J}^{(0)}(k^+) = \delta(k^+), \quad \mathcal{J}^{(-2)}(k^+_j) = \frac{\delta(k^+_i) - \delta(k^+_2)}{k^+_2 - k^+_1}, \quad \mathcal{J}^{(-4)}(k^+_j) = 4\pi\alpha_s(\mu_i) \left[ \frac{\delta(k^+_i)}{(k^+_2)(k^+_3)} + \frac{\delta(k^+_2)}{(k^+_1)(k^+_3)} + \frac{\delta(k^+_3)}{(k^+_1)(k^+_2)} - \pi^2\delta(k^+_i)\delta(k^+_2)\delta(k^+_3) \right].
\]

At one-loop order, we have

\[
\mathcal{J}^{(0)}(\omega, k^+, \mu) = \left\{ \delta(k^+) \left[ 1 + \frac{\alpha_s(\mu)C_F}{4\pi} \left( 2\ln^2 \frac{\omega_{pX}}{\mu^2} - 3\ln \frac{\omega_{pX}}{\mu^2} + 7 - \pi^2 \right) \right] + \frac{\alpha_s(\mu)C_F}{4\pi} \left( 4\ln \frac{\omega_{pX}}{k^+} + 3\ln \frac{\omega_{pX}}{\mu^2} - 7 - \pi^2 \right) \right\} \theta(p^+_X - k^+) \theta(k^+),
\]

where \( \omega = \bar{n} \cdot p \) is the large partonic momentum.

For convenience we define

\[
F(p^+, p^-) = \int_0^{p^+} dk^+ \mathcal{J}^{(0)}(p^-, k^+, \mu) f^{(0)}(k^+ + \bar{\Lambda} - p^+, \mu)
\]

\[
+ \frac{1}{2m_B} f^{(2)}_0(\bar{\Lambda} - p^+) - \frac{\lambda_1 + 3\lambda_2}{2m_B} f^{(0)}(\bar{\Lambda} - p^+),
\]

\[
F_{1,2}(p^+) = f^{(2)}_{1,2}(\bar{\Lambda} - p^+),
\]

where a prime denotes a derivative, as well as

\[
F_{3,4}(p^+) = \int dk^+_1 dk^+_2 \left[ \delta(k^+_1) - \delta(k^+_2) \right] f^{(4)}_{3,4}(k^+_j + \bar{\Lambda} - p^+),
\]

\[
F_{5,6}(p^+) = \int dk^+_1 dk^+_2 dk^+_3 \left[ \delta(k^+_1) + \delta(k^+_2) + \delta(k^+_3) \right] f^{(6)}_{5,6}(k^+_j + \bar{\Lambda} - p^+).
\]

If we use the tree-level expression for \( \mathcal{J}^{(0)} \), then \( F(p^+, p^-) = F(p^+) \) is a function of \( p^+ \) only. Then, for \( B \to X_s\gamma \), the rate \( d\Gamma^s/dx_H^\gamma \) in the endpoint region is \cite{32} \footnote{This includes \( \mathcal{O}_7 - \mathcal{O}_7 \) and \( \mathcal{O}_7 - \mathcal{O}_2 \) contributions only. }

\[
\frac{1}{\Gamma_s^0} \frac{d\Gamma^s}{dx_H^\gamma} \bigg|_{x_H^\gamma > x_H} = m_B(C^{(i)})^2 \left[ 1 - 3(1 - x_H^\gamma) \right] F(m_B(1 - x_H^\gamma), m_B)
\]

\[
+ \left[ m_B(1 - x_H^\gamma) - \bar{\Lambda} \right] F(m_B(1 - x_H^\gamma)) + F_2(m_B(1 - x_H^\gamma))
\]

\[
- F_3(m_B(1 - x_H^\gamma)) + F_4(m_B(1 - x_H^\gamma)) - 8\pi\alpha_s(\mu_i) F_5^\gamma(m_B(1 - x_H^\gamma)),
\]

\footnote{This includes \( \mathcal{O}_7 - \mathcal{O}_7 \) and \( \mathcal{O}_7 - \mathcal{O}_2 \) contributions only. In Ref. \cite{38} subleading corrections from \( \mathcal{O}_7 - \mathcal{O}_8 \) are studied and estimated to contribute between \(-0.3\%\) and \(-3\%\) to the total flavour-averaged decay rate. We do not consider such corrections in this work.}
where $1 - x_H^c \sim \lambda^2$ and

$$C^{(i)} = 1 + \Delta_\gamma(m_b, \varrho) - \frac{\alpha_s(m_b)C_F}{4\pi} \left\{ \frac{\pi^2}{12} + 6 \right\},$$

$$\Delta_\gamma(m_b, \varrho) = \frac{1}{C_7^{\text{eff}(0)}(m_b)} \left\{ \frac{\alpha_s}{4\pi} C_7^{\text{eff}(1)}(m_b) + \sum_k C_k^{\text{eff}(0)}(m_b)r_k(\varrho) \right\}. \quad (12)$$

The triply differential decay rate for $B \to X_u \ell \bar{\nu}$ at NLO [32] is obtained by substituting the $W_i^{\ell\nu}$ listed in Appendix B into Eq. (5). At tree level, this becomes

$$\frac{1}{\Gamma_0} \frac{d^3\Gamma^u}{dx_H d\bar{y}_H du_H} = 6(1 - u_H)(x_H + \bar{y}_H - 1) \left\{ 2m_B(2 - x_H - \bar{y}_H - u_H)F(m_B u_H) - \frac{1}{\bar{y}_H - u_H} \left( \bar{y}_H^2 - (2 - x_H)\bar{y}_H + 2(1 - x_H) - u_H(2 - x_H - u_H) \right) F_1(m_B u_H) + \frac{2}{\bar{y}_H(\bar{y}_H - u_H)} \left( \bar{y}_H^3 - (2 - x_H)\bar{y}_H^2 - (4 - u_H)(x_H + u_H)\bar{y}_H + 2(x_H + 2\bar{y}_H + u_H - 1) \right) F_2(m_B u_H) + \frac{2}{\bar{y}_H}(x_H + \bar{y}_H + u_H - 2) F_3(m_B u_H) - \frac{2}{\bar{y}_H(\bar{y}_H - u_H)} \left( \bar{y}_H^2 - (2 - x_H)\bar{y}_H + 2(1 - x_H) - u_H(2 - x_H - u_H) \right) F_4(m_B u_H) - \frac{4}{\bar{y}_H(\bar{y}_H - u_H)} (1 - \bar{y}_H)(x_H + \bar{y}_H - 1) 4\pi \alpha_s(\mu_i)F_5^{\mu}(m_B u_H) + \frac{4}{\bar{y}_H(\bar{y}_H - u_H)} (1 - u_H)(1 - x_H - u_H) 4\pi \alpha_s(\mu_i)F_6^{\mu}(m_B u_H) \right\}. \quad (13)$$

Note that we can use the relation [30]

$$F_1(m_B u_H) = 2(\bar{A} - m_B u_H)F(m_B u_H) + \mathcal{O}(\lambda^4) \quad (14)$$

to eliminate $F_1(m_B u_H)$, as was done in Eq. (11).

The triply differential decay rate for $B \to X_s \ell^+ \ell^-$ was calculated in Refs. [24, 36]. The $W_i^{\ell\nu}$ appearing in Eq. (5) are also listed in Appendix B.

III. $|V_{ub}|$ AT NLO

A. Relations between $B \to X_u \ell \bar{\nu}$ and $B \to X_s \gamma$

Consider first the process $B \to X_u \ell \bar{\nu}$. We wish to isolate or eliminate the subleading shape functions that appear in the rates. In the following, we shall work at tree level. Inspection of Eqs. (B2) and (B9) shows that the shape functions appear in the hadronic structure functions $W_1$ to $W_3$ in only two combinations, namely

$$m_B F_I = m_B F + \frac{1}{2} F_1 - F_2 - \frac{1}{\bar{y}_H} (F_3 - F_4 + \frac{8\pi \alpha_s(\mu_i)F_5^{\mu}}{\bar{y}_H}),$$

$$m_B F_{II} = F_1 - \frac{2(\bar{y}_H(2 - u_H) - 1)}{\bar{y}_H(1 - u_H)} F_2 + \frac{2}{\bar{y}_H} (F_4 - 4\pi \alpha_s(\mu_i)F_5^{\mu} - 4\pi \alpha_s(\mu_i)F_6^{\mu}), \quad (15)$$
where we have suppressed the argument $m_B u_H$. Specifically,

$$
W_1 = \frac{1}{4} F_I ,
$$

$$
W_2 = \frac{1 - u_H}{\bar{y}_H - u_H} F_I - \frac{(1 - u_H)^2}{(\bar{y}_H - u_H)^2} F_{II} ,
$$

$$
W_3 = \frac{1}{2m_B (\bar{y}_H - u_H)} F_I .
$$

(16)

Nevertheless, taking integrals of the form

$$
\int_{u_H}^{1} d\bar{y}_H \int_{1 - u_H}^{1 - \bar{y}_H} dx_H K^u(x_H, \bar{y}_H, u_H) \frac{d^3 \Gamma^u}{dx_H d\bar{y}_H d u_H} ,
$$

with suitable choices of the weight function $K^u(x_H, \bar{y}_H, u_H)$, we can isolate the following four linearly independent combinations of the $F_i$:

$$
(4 - 2u_H)m_B F + F_1 ,
$$

$$
(1 - u_H)m_B F + F_2 ,
$$

$$
F_3 - F_4 + 8\pi \alpha_s(\mu_i) F_5^u ,
$$

$$
m_B F - \frac{1}{2} F_3 - \frac{1}{2} F_4 + 4\pi \alpha_s(\mu_i) F_6^u .
$$

(18)

(Recall that we can apply Eq. (14) so that the first combination involves only the leading-order shape function.) Here, the treatment of the $u_H$ dependence in the rate requires care. Expanding Eq. (13) in $u_H \sim \lambda^2$ when obtaining the weight function will typically result in excessively large coefficients in the $u_H F_{1-6}(m_B u_H)$ terms (which are formally of order $\lambda^4$). For example, choosing $K^u(x_H, \bar{y}_H) = -21 x_H + 21 \bar{y}_H + 45 x_H \bar{y}_H - \frac{25}{2} y_H^2$, we obtain

$$
\frac{1}{\Gamma_0} \int \int dx_H d\bar{y}_H K^u(x_H, \bar{y}_H) \frac{d^3 \Gamma^u}{dx_H d\bar{y}_H d u_H} = (1 - 7u_H)m_B F(m_B u_H) + \frac{1}{4} F_1(m_B u_H) + O(\lambda^4) ,
$$

(19)

so this eliminates all but the leading-order shape function up to $O(\lambda^4)$ corrections. However, we then have the additional contributions

$$
\frac{5}{4} u_H F_1(m_B u_H) - \frac{49}{2} u_H F_2(m_B u_H) - \frac{109}{4} u_H F_3(m_B u_H) + \frac{57}{4} u_H F_4(m_B u_H)
$$

$$
- \frac{83}{2} u_H \times 4\pi \alpha_s(\mu_i) F_5^u(m_B u_H) + 13 u_H \times 4\pi \alpha_s(\mu_i) F_6^u(m_B u_H) .
$$

(20)

For this reason, when calculating $K^u$, we keep the full dependence on $u_H$ in the rate, rather than dropping terms that are formally subleading in a strict SCET expansion in $u_H/\bar{y}_H \sim \lambda^2$. (The analysis of $m_X$-cut effects in $B \rightarrow X_s \ell^+ \ell^-$ [24, 25] also retained the full $u_H$ dependence, since doing so facilitates making contact with the total rate in the local OPE [39, 40, 41].) Thus, subleading shape functions are eliminated to all orders in $u_H$, and the issue is resolved. One straightforward method for obtaining $K^u(x_H, \bar{y}_H, u_H)$ is then to take different moments of the rate with respect to $x_H$ and $\bar{y}_H$, and solve the resulting set of linear equations in the $F_i$. 

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In Eq. (17), we consider the case where a cut is imposed on $p_X^+$, i.e. $p_X^+ < m_B^2/m_B$. Different or additional cuts will change the limits of integration, calling for different weight functions. Table I lists several examples of $K^u$s that isolate the combination $m_B F + F_1/(4 - 2u_H)$, while Tables II and III give examples that result in (18b) and (18c) respectively.

\begin{align*}
(1) \quad K_1^u &= \frac{5}{9} \frac{1}{\left(1-u_H\right)^8} \left[ 10(7 - u_H)(1 - u_H)(4 + 3u_H)y_H \\
&\quad - (454 + 247u_H - 71u_H^2)x_H y_H - 4(1 - u_H)(109 - 4u_H)y_H^2 \\
&\quad + 105(7 - u_H)x_H y_H^2 \right] \\
(2) \quad K_1^u &= \frac{5}{32} \frac{1}{\left(1-u_H\right)^8} \left[ -10(7 - u_H)(1 - u_H)(34 - 27u_H)y_H \\
&\quad + 2(1 - u_H)(2759 - 449u_H)x_H y_H - 525(7 - u_H)x_H^2 y_H \\
&\quad + 2(1 - u_H)(341 - 131u_H)y_H^2 \right] \\
(3) \quad K_1^u &= \frac{15}{41} \frac{1}{\left(1-u_H\right)^8} \left[ -2(1 - u_H)^2(288 - 29u_H)y_H \\
&\quad + (26 - 1793u_H + 157u_H^2)x_H y_H - 10(109 - 4u_H)x_H^2 y_H \\
&\quad + (341 - 131u_H)x_H y_H^2 \right]
\end{align*}

### TABLE I

| Example | Expression |
|---------|------------|
| (A)     | $K_1^u = \frac{5}{9} \frac{1}{\left(1-u_H\right)^8} \left[ -2(1 - u_H)(7 - 15u_H)y_H - (34 + 71u_H)x_H y_H \\
&\quad - 16(1 - u_H)y_H^2 + 105x_H y_H^2 \right]$ |
| (B)     | $K_1^u = \frac{5}{32} \frac{1}{\left(1-u_H\right)^8} \left[ -2(1 - u_H)(266 - 135u_H)y_H + 898(1 - u_H)x_H y_H \\
&\quad - 525x_H^2 y_H + 262(1 - u_H)y_H^2 \right]$ |
| (C)     | $K_1^u = \frac{15}{41} \frac{1}{\left(1-u_H\right)^8} \left[ -58(1 - u_H)^2 y_H + (26 - 157u_H)x_H y_H \\
&\quad - 40x_H^2 y_H + 131x_H y_H^2 \right]$ |

### TABLE II

| Example | Expression |
|---------|------------|
| (A)     | $K_1^u = \frac{5}{9} \frac{1}{\left(1-u_H\right)^8} \left[ -2(1 - u_H)(7 - 15u_H)y_H - (34 + 71u_H)x_H y_H \\
&\quad - 16(1 - u_H)y_H^2 + 105x_H y_H^2 \right]$ |
| (B)     | $K_1^u = \frac{5}{32} \frac{1}{\left(1-u_H\right)^8} \left[ -2(1 - u_H)(266 - 135u_H)y_H + 898(1 - u_H)x_H y_H \\
&\quad - 525x_H^2 y_H + 262(1 - u_H)y_H^2 \right]$ |
| (C)     | $K_1^u = \frac{15}{41} \frac{1}{\left(1-u_H\right)^8} \left[ -58(1 - u_H)^2 y_H + (26 - 157u_H)x_H y_H \\
&\quad - 40x_H^2 y_H + 131x_H y_H^2 \right]$ |

TABLE I: Some choices of $K^u(x_H, y_H, u_H)$ for which the weighted integral Eq. (17) equals $m_B F + F_1/(4 - 2u_H)$.  

TABLE II: Some choices of $K^u(x_H, y_H, u_H)$ for which the weighted integral Eq. (17) equals $(1 - u_H)m_B F + F_2$. 

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TABLE III: Some choices of $K_u(x_H, \bar{y}_H, u_H)$ for which the weighted integral Eq. (17) equals $F_3 - F_4 + 2\tilde{F}_5^u$. 

Now, the subleading shape functions $F_{5,6}$ depend upon the light-quark flavour (see Appendix A). We indicate this difference between the $F_{5,6}'s$ appearing in $B \to X_u\ell\bar{\nu}$ and $B \to X_s\gamma$ by using the superscripts ‘u’ and ‘s’. In order to cancel the $F_s^u$ contribution to the latter decay,\(^2\) we can use approximate SU(3) flavour symmetry, namely the fact that

$$\frac{F_u^s - F_s^s}{F_s^s} \sim \frac{m_s}{\Lambda_{QCD}} \quad \text{(21)}$$

is suppressed. This enables us to relate the semileptonic process to the radiative process and thereby derive an expression for $\Gamma_0^u$, or equivalently $|V_{ub}|$, to subleading order. We can write

$$\frac{1}{\Gamma_0^u} \int \int \left[ K_u^u - K_{\|}^u + \rho K_{\perp}^u \right] \frac{d^3\Gamma^u}{dx_Hd\bar{y}_Hdu_H} \, dx_Hd\bar{y}_H = m_B F(m_Bu_H) - \frac{1}{2} \left[ F_1(m_Bu_H) - 2F_2(m_Bu_H) \right] - \left[ F_3(m_Bu_H) - F_4(m_Bu_H) + 2\tilde{F}_5^u(m_Bu_H) \right], \tag{22}$$

\(^2\) The authors of Refs. [33, 34] have used model-dependent arguments to estimate that the effects of $f_{5,6}$, when integrated over a sufficiently large region, are comparatively small ($\sim 5\%$), but that they may cause large corrections in the $d\Gamma/dp_X^+$ spectrum for $p_X^+ \leq 0.5$ GeV. We avoid any need to consider the reliability of these numerics by simply eliminating $f_{5,6}$, along with the other tree-level shape functions.
where

$$\rho(u_H) = \frac{(2 - u_H)(u_H - \frac{\Lambda}{2m_B})}{(1 - u_H + \frac{\Lambda}{2m_B})} \quad (23)$$

and $\tilde{F}_5^u = 4\pi\alpha_s(\mu_i)F_5^u$. $K_1^u$, $K_2^u$ and $K_3^u$ are any weight functions that give the combinations $m_B F + F_1/(4-2u_H)$, $(1-u_H)m_B F + F_2$ and $F_3 - F_2 + 2\tilde{F}_5^u$ respectively (examples of which are presented in Tables I, II and III). The shape functions in Eq. (22) appear in the same linear combination as in the rate $d\Gamma^s/du_H$. Hence, at NLO we obtain

$$\frac{1}{\Gamma_0^u} \int \left[ K_2^u - K_3^u + \rho K_1^u \right] \frac{d^3\Gamma^u}{dx_H d\bar{y}_H du_H} dx_H d\bar{y}_H = -\frac{1}{(1-u_H)^3} \frac{1}{\Gamma_0^s} \frac{d\Gamma^s}{du_H}. \quad (24)$$

More generally, we can construct $K^u$ such that

$$\hat{M}^u \equiv \frac{1}{\Gamma_0^u} M^u \equiv \frac{1}{\Gamma_0^u} \int \int K^u(x_H, \bar{y}_H, u_H) \frac{d^3\Gamma^u}{dx_H d\bar{y}_H du_H} dx_H d\bar{y}_H \quad (25)$$

$$= m_B F(m_B u_H) + \kappa_1^u(u_H) F_1(m_B u_H)$$

$$+ \kappa_2^u(u_H) [F_2(m_B u_H) - F_3(m_B u_H) + F_4(m_B u_H) - 2\tilde{F}_5^u(m_B u_H)].$$

For example, we can use

$$K^u = K_N^u - \kappa_2^uK_3^u, \quad (26)$$

where $K_N^u$ is a weight function that gives the linear combination $m_B F + \kappa_1^u F_1 + \kappa_2^u F_2$, examples of which are given in Table IV in Appendix C (with the corresponding values of $\kappa_1^u$ and $\kappa_2^u$ shown there). We can also use

$$K^u = \beta K_1^u + \frac{1 - \beta}{1 - u_H} (K_2^u - K_3^u), \quad (27)$$

with $\beta$ an arbitrary real number (in which case $\kappa_1^u = \beta/(4-2u_H)$ and $\kappa_2^u = (1-\beta)/(1-u_H)$).

For any such $K^u$, we have

$$\hat{M}^u + \kappa_2^u (1-u_H)^{-3} \hat{M}^s = \left\{ (1 - \kappa_2^u) + \left( \frac{\Lambda}{m_B} - u_H \right) \left( 2\kappa_1^u + \kappa_2^u \right) \right\} m_B F(m_B u_H) \quad (28)$$

$$+ \mathcal{O}(\alpha_s, \lambda^4),$$

where $\hat{M}^s = (1/\Gamma_0^s) M^s = (1/\Gamma_0^s)(d\Gamma^s/du_H)$, i.e. combining $\hat{M}^u$ and $\hat{M}^s$ in this way gives an expression dependent only on the leading-order shape function. Taking the ratio of two such expressions (two choices of $K^u$) at $u_H \neq 0$ then provides us with a relation independent of both leading and subleading shape functions. We shall use the superscripts $(i)$ and $(ii)$ when we need to distinguish between quantities in the two expressions. We then obtain

$$\frac{\Gamma_0^u}{\Gamma_0^s} = -\frac{\left[ b_0^{(ii)} M^u(i) - b_0^{(i)} M^u(i) \right]}{\left[ b_0^{(ii)} \kappa_2^u(i) - b_0^{(i)} \kappa_2^u(i) \right]} (1-u_H)^{-3} M^s, \quad (29)$$

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where
\[ b_0 = (1 - \kappa_u^u) + \left( \frac{\bar{\Lambda}}{m_B} - u_H \right) \left( 2\kappa_1^u + \kappa_2^u \right). \] (30)

Since the right-hand side of Eq. (29) is measurable, it enables an experimental determination of the CKM ratio on the left-hand side. Additionally, the factor $|V_{tb}V_{ts}^\ast|$ in this ratio can be eliminated by normalizing the photon spectrum by the total $B \to X_s\gamma$ rate, which is given in a local OPE.

There will be loop and power ($\lambda^4$-suppressed) corrections to the rates and hence also to Eq. (29). While these are not fully known, one can show that the corrections to Eq. (29) are proportional to
\[ b_0^{(i)} \kappa_2^{u(ii)} - b_0^{(ii)} \kappa_2^{u(i)} + \cdots \] (31)
(multiplied by $\alpha_s$ or $\lambda^4$). This needs to be taken into account when selecting $\{K_{u}^{(i)}, K_{u}^{(ii)}\}$: one should avoid pairs of weight functions that result in Eq. (31) being excessively large, lest parametrically suppressed terms acquire excessively large coefficients. For example, one appropriate choice is to use Eq. (27) for both $K_{u}^{(i)}$, with $\beta^{(i)} = 1$ and $\beta^{(ii)} = 0$, after which the magnitude of Eq. (31) is less than $1/6$ for $0 < u_H < m_D^2/m_B^2$.

**B. Relations involving $B \to X_s\ell^+\ell^-$**

We can also try to isolate shape functions in the process $B \to X_s\ell^+\ell^-$ by taking integrals of the form
\[ \int_{\tilde{y}_{\min}}^{\tilde{y}_{\max}} d\tilde{y}_H \int_{1-\tilde{y}_H}^{1-u_H} dx_H K^{\ell\ell}(x_H, \tilde{y}_H, u_H) \frac{d^3\Gamma^{\ell\ell}}{dx_H d\tilde{y}_H du_H}, \] (32)
where
\[ \tilde{y}_{\min(max)} = 1 - \frac{y_{\max(min)}}{1-u_H}. \] (33)

Here, $y_H = q^2/m_B^2$ and the low-$q^2$ region corresponds to $1 \text{GeV}^2 \leq q^2 \leq 6 \text{GeV}^2$. However, determining $K^{\ell\ell}(x_H, \tilde{y}_H, u_H)$ in the straightforward manner described above proves to be problematic in practice. Therefore, we resort to another method, which is based on the following observation. Under the transformation $x_H \to x'_H = 2 - u_H - \tilde{y}_H - x_H$, we find that
\[ \int_{1-\tilde{y}_H}^{1-u_H} dx_H = \int_{1-\tilde{y}_H}^{1-u_H} dx'_H \quad \text{and} \]
\[ (1 - x_H - u_H) \leftrightarrow (x_H + \tilde{y}_H - 1), \]
\[ (2x_H + u_H + \tilde{y}_H - 2) \leftrightarrow -(2x_H + u_H + \tilde{y}_H - 2). \]

This symmetry or antisymmetry can be exploited to obtain $K^{\ell\ell}$. For example, if $K^{\ell\ell}$ changes sign under the transformation, then we can see from the triply differential rate, Eq. (5), that integration over $x_H$ eliminates the $W_1$ and $W_2$ terms, whereas the $W_3$ term remains. Now,
Eq. (B10) shows that \( F_3, F_4 \) and \( F_5^\ast \) occur in \( W_3 \) in the same linear combination as in the \( B \to X_s \gamma \) rate.

This still leaves the integration over \( \bar{y}_H \), and if we choose \( K^{\ell\ell}(x_H, \bar{y}_H, u_H) = (2x_H + u_H + \bar{y}_H - 2) \tilde{K}^{\ell\ell}(\bar{y}_H, u_H) \), where \( \tilde{K}^{\ell\ell}(y, u_H) \) satisfies

\[
\int_{y_{\min}}^{y_{\max}} d\bar{y}_H(\bar{y}_H - u_H)^3 \frac{1}{\bar{y}_H} (2\text{Re}[C_{10a} C_{1a}^*] + \text{Re}[C_{10}\phi] (1 - \bar{y}_H^2)) \tilde{K}^{\ell\ell}(\bar{y}_H, u_H) = 0, \tag{34}
\]

then all of the subleading shape functions in Eq. (32) appear in the same combination as in the \( B \to X_s \gamma \) rate, which can thus be used to eliminate these functions. Table V in Appendix C shows several examples of \( K^{\ell\ell} \) of this form. We observe that \( z = \cos \theta = (2x_H + u_H + \bar{y}_H - 2)/(\bar{y}_H - u_H) \), where \( \theta \) is the angle between the \( B \) and \( \ell^+ \) in the center-of-mass frame of the \( \ell^+ \ell^- \). This means that a choice of \( K^{\ell\ell} \propto (2x_H + u_H + \bar{y}_H - 2) \) is equivalent to taking moments of the forward-backward asymmetry,

\[
\frac{d^2 A_{FB}}{dy_H du_H} = \int_{-1}^{1} dz \frac{\text{sign}(z)}{\Gamma_0} \frac{d^3 \Gamma}{dy_H du_H dz} = \frac{3}{2\Gamma_0} \int_{-1}^{1} dz z \frac{d^3 \Gamma}{dy_H du_H dz}. \tag{35}
\]

Note also that \( C_{9a} \) is a function of \( q^2 \), and hence of \( \bar{y}_H \) (see Appendix B), but in the low-\( q^2 \) region \( |C_{9a}| \) varies by less than \( \pm 1\% \) and we take it to be constant. There is no problem taking into account the exact dependence, but integrals over regions of \( \bar{y}_H \) must then be performed numerically.

Let \( \Gamma_0 M^s = d\Gamma^s / du_H \), and let \( M^u = \Gamma_0 \tilde{M}^u \) and \( \Gamma_0 \tilde{M}^{\ell\ell} \) denote the integrals (17) and (32) respectively, with weight functions from Tables I and V. Then we obtain

\[
\Gamma_0^u = \frac{1 + \kappa_3^{\ell\ell}}{1 + 2(\frac{\Delta}{m_B} - u)\kappa_1^u} \frac{M^u}{\tilde{M}^{\ell\ell} - \kappa_3(1 - u_H)^{3/2} M^s}, \tag{36}
\]

where \( \kappa_1^u (\kappa_3^{\ell\ell}) \) is the coefficient of \( F_1 (F_3) \) in \( \tilde{M}^u (\tilde{M}^{\ell\ell}) \).

More generally, by the same methods, we can find \( K^u \) and \( K^{\ell\ell} \) such that

\[
\tilde{M}^u \equiv \frac{1}{\Gamma_0^u} M^u \equiv \frac{1}{\Gamma_0^u} \int K^u(x_H, \bar{y}_H, u_H) \frac{d^3 \Gamma^u}{dx_H dy_H du_H} dx_H d\bar{y}_H \
= m_B F(m_B u_H) + \kappa_1^u(u_H) F_1(m_B u_H) + \kappa_2^u(u_H) F_2(m_B u_H),
\]

\[
\tilde{M}^s \equiv \frac{1}{\Gamma_0^s} M^s \equiv \frac{1}{\Gamma_0^s} \frac{d\Gamma^s}{du_H} \
= -(1 - u_H)^3 \left\{ m_B F(m_B u_H) - \frac{1}{2} \left[ F_1(m_B u_H) - 2 F_2(m_B u_H) \right] \right\} 
- \left[ F_3(m_B u_H) - F_4(m_B u_H) + 2 \tilde{F}_5^s(m_B u_H) \right] ,
\]

\[
\tilde{M}^{\ell\ell} \equiv \frac{1}{\Gamma_0^{\ell\ell}} M^{\ell\ell} \equiv \frac{1}{\Gamma_0^{\ell\ell}} \int K^{\ell\ell}(x_H, \bar{y}_H, u_H) \frac{d^3 \Gamma^{\ell\ell}}{dx_H dy_H du_H} dx_H d\bar{y}_H \
= m_B F(m_B u_H) + \frac{1}{2} \kappa_2^{\ell\ell}(u_H) \left[ F_1(m_B u_H) - 2 F_2(m_B u_H) \right] 
+ \kappa_3^{\ell\ell}(u_H) \left[ F_3(m_B u_H) - F_4(m_B u_H) + 2 \tilde{F}_5^s(m_B u_H) \right] ,
\]

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where \( \tilde{F}_5^s = 4\pi\alpha_s(\mu_i)F_5^s \). Tables IV and VI show (further) examples of such weight functions, along with the corresponding values of the coefficients \( \kappa_{1,2}^u \) and \( \kappa_{2,3}^{\ell\ell} \). Then

\[
\left[ \kappa_2^{\ell\ell} - \kappa_3^{\ell\ell} \right] \check{M}^u + \kappa_2^u \check{M}^{\ell\ell} - \kappa_2^{\ell\ell} \kappa_3^u (1 - u_H)^{-3} \check{M}^s
\]

\[
= \left\{ \kappa_2^{\ell\ell} + \kappa_2^u + \kappa_2^{\ell\ell} \kappa_3^u \right\} + \left( \frac{\bar{\Lambda}}{m_B} - u_H \right) \left( \kappa_2^{\ell\ell} - \kappa_3^u \right) (2\kappa_1^u + \kappa_2^u) \left\{ m_B F (m_B u_H) + \mathcal{O}(\alpha_s, \lambda^4) \right\},
\]

so in this case we have a combination of \( \check{M}^u \), \( \check{M}^s \) and \( \check{M}^{\ell\ell} \) that is dependent only on the leading-order shape function. Taking the ratio of two such expressions (two choices of \( \{K^u, K^{\ell\ell}\} \), denoted by superscripts \( (i) \) and \( (ii) \) as previously) at \( u_H \neq 0 \) then provides us with another relation independent of both leading and subleading shape functions.

Specifically, let

\[
\check{\mathcal{M}}^u = \frac{1}{\Gamma_0^u} M^u = \left\{ \begin{array}{ll} \kappa_2^{\ell\ell} - \kappa_3^{\ell\ell} [\check{M}^u], & \text{if } \kappa_2^u \neq 0 \\ \check{M}^u, & \text{if } \kappa_2^u = 0 \end{array} \right.,
\]

\[
\check{\mathcal{M}}^{\ell\ell} = \frac{1}{\Gamma_0^{\ell\ell}} M^{\ell\ell} = \left\{ \kappa_2^{\ell\ell} [\check{M}^{\ell\ell}], & \text{if } \kappa_2^{\ell\ell} \neq \kappa_3^{\ell\ell} \\ \check{M}^{\ell\ell}, & \text{if } \kappa_2^{\ell\ell} = \kappa_3^{\ell\ell} \right..
\]

\[
\check{\mathcal{M}}^s = \frac{1}{\Gamma_0^s} M^s = \left\{ \begin{array}{ll} -\kappa_2^{\ell\ell} \kappa_3^u (1 - u_H)^{-3} [\check{M}^s], & \text{if } \kappa_2^{\ell\ell} \neq \kappa_3^u \\ -\kappa_3^u (1 - u_H)^{-3} [\check{M}^s], & \text{if } \kappa_2^{\ell\ell} = \kappa_3^u \end{array} \right.,
\]

and

\[
c_0 = \left\{ \begin{array}{ll} (\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell} + \kappa_2^u + \kappa_2^{\ell\ell} \kappa_3^u) + \left( \frac{\bar{\Lambda}}{m_B} - u_H \right) \left( \kappa_2^{\ell\ell} - \kappa_3^u \right) (2\kappa_1^u + \kappa_2^u), & \text{if } \kappa_2^u \neq 0 \text{ and } \kappa_2^{\ell\ell} \neq \kappa_3^u \\ 1 + 2 \left( \frac{\bar{\Lambda}}{m_B} - u_H \right) \kappa_1^u, & \text{if } \kappa_2^u = 0 \\ (1 + \kappa_3^{\ell\ell}), & \text{if } \kappa_2^{\ell\ell} = \kappa_3^{\ell\ell} \end{array} \right..
\]

We find that

\[
\frac{\Gamma_0^u}{\Gamma_0^{\ell\ell}} = - \left[ \frac{c_0^{(ii)} \mathcal{M}^{u(ii)} - c_0^{(i)} \mathcal{M}^{u(ii)}}{c_0^{(ii)} (\mathcal{M}^{s(ii)} + r \mathcal{M}^{\ell\ell(ii)}) - c_0^{(i)} (\mathcal{M}^{s(ii)} + r \mathcal{M}^{\ell\ell(ii)})} \right],
\]

where \( r = \Gamma_0^s / \Gamma_0^{\ell\ell} \), or

\[
\frac{\Gamma_0^u}{\Gamma_0^s} = - \left[ \frac{c_0^{(ii)} \mathcal{M}^{u(ii)} - c_0^{(i)} \mathcal{M}^{u(ii)}}{c_0^{(ii)} (\mathcal{M}^{s(ii)} + \mathcal{M}^{\ell\ell(ii)}) - c_0^{(i)} (\mathcal{M}^{s(ii)} + \mathcal{M}^{\ell\ell(ii)})} \right].
\]

In the special case where \( \kappa_2^u = 0 \) and \( \kappa_2^{\ell\ell} = \kappa_3^{\ell\ell} \), Eq. (42) reduces to Eq. (36).

The loop and power (\( \lambda^4 \)-suppressed) corrections to Eq. (41) can be shown to be proportional to

\[
- \frac{c_0^{(ii)} [\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell}]^{(ii)}}{c_0^{(ii)} [\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell}]^{(i)} - c_0^{(i)} [\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell}]^{(i)}} + \frac{c_0^{(i)} [\kappa_2^u (1 + \kappa_3^{\ell\ell})]^{(ii)}}{c_0^{(i)} [\kappa_2^u (1 + \kappa_3^{\ell\ell})]^{(i)} - c_0^{(i)} [\kappa_2^u (1 + \kappa_3^{\ell\ell})]^{(i)}} + \ldots,
\]
where \( \tilde{c}_0 = (\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell} + \kappa_2^{u} + \kappa_2^{u} \kappa_3^{\ell\ell}) + (\bar{\Lambda}/m_B - u_H) (\kappa_2^{\ell\ell} - \kappa_3^{\ell\ell})(2\kappa_1^{u} + \kappa_2^{u}). \) When selecting \( \{K^{u(i)}_u, K^{\ell\ell(ii)}_\ell\ell\} \), \( \{K^{u(ii)}_u, K^{\ell\ell(ii)}_\ell\ell\} \), one should avoid those sets of weight functions that result in Eq. (43) being excessively large. The following combinations of weight functions are suitable choices:

\[
K^{u(i)}_u = (1), (2) \text{ or } (3) \quad \text{[Table I]} \quad \text{and} \quad K^{\ell\ell(ii)}_\ell\ell = (7), (8) \text{ or } (9) \quad \text{[Table V]};
\]
\[
K^{u(ii)}_u = (4), (5) \text{ or } (6) \quad \text{[Table IV]} \quad \text{and} \quad K^{\ell\ell(ii)}_\ell\ell = (7), (8) \text{ or } (9) \quad \text{[Table V]}.
\]

C. Perturbative Corrections

Let us now consider the feasibility of incorporating perturbative corrections in our relations. In Ref. [32], the complete set of subleading corrections (to all orders in \( \alpha_s \)) for the triply differential spectrum of \( B \to X_u \ell\bar{\nu} \) was derived. It was shown that prohibitively many new shape functions appear at order \( \alpha_s \Lambda_{\text{QCD}}/m_b \), and hence it is not phenomenologically viable to work to that order.\(^3\) However, one may choose to work to order \( \alpha_s \lambda^0, \alpha_s^0 \lambda^2 \), by including perturbative corrections to just the leading-power terms. Recall that there are two perturbative scales, \( \mu_b \sim m_b \) (hard) and \( \mu_i \sim \sqrt{m_b \Lambda_{\text{QCD}}} \) (jet). It is straightforward to take into account the relevant hard corrections. Including the effect of corrections to the jet function \( J^{(0)} \), which is convoluted with the shape function \( f^{(0)} \), is more involved: one has to “invert” a distribution (see Eq. (8)). An implementation akin to Refs. [26, 27, 28, 29] is left for future work. Nevertheless, before this is done, we can still use the less direct approach mentioned in the introduction, using two instances of Eq. (28) or (38), with appropriately modified right-hand sides. For example, one can extract the leading-order shape function from the analogue of Eq. (38), with \( K^{\ell\ell} \) from Table V, and substitute this function into a second choice, with \( K^u \) from Table I. Finally, we note that the extent to which Eq. (29) or (42) varies with respect to \( u_H \) or different combinations of the \( K^u \)s and \( K^{\ell\ell} \)s will provide a measure of the effect of \( \alpha_s \) and \( \lambda^4 \) corrections.

IV. CONCLUSION

In this paper, we have established a method for obtaining \( |V_{ub}|/|V_{tb}V_{ts}^*| \) that includes \( \mathcal{O}(\Lambda_{\text{QCD}}/m_b) \) corrections in a model-independent way. Our approach relies upon a class of relations between the inclusive decays \( B \to X_u \ell\bar{\nu} \) and \( B \to X_s \gamma \) that are valid including the first-order power corrections (see Eqs. (24) and (29)). Alternatively, one can use a separate class of relations involving \( B \to X_u \ell^+ \ell^- \) (see Eqs. (36) and (42)). Experimentally required cuts make shape-function effects important in these processes. Their differential decay spectra in the shape function region have previously been derived to subleading order with the help of the Soft-Collinear Effective Theory. These rates involve a number of non-perturbative but universal shape functions in different linear combinations. We are able to eliminate these sources of hadronic uncertainty by taking suitable weighted integrals of the triply differential rates. Hence, our results incorporate NLO power corrections while

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\(^3\) Unless these shape functions appear in the rates in only a much smaller number of linear combinations.
avoiding model dependence. There are many possible weight functions (see e.g. Eqs. (26) and (27)); different choices provide a consistency check on the determination of \(|V_{ub}|\).

Acknowledgments

I wish to thank Iain Stewart for helpful discussions and for comments on the manuscript. I am also grateful for conversations with Antonio Limosani, Frank Tackmann and Mark Wise.

APPENDIX A: SHAPE FUNCTIONS

The leading-order shape function is
\[ f^{(0)}(\ell^+) = \frac{1}{2} \langle B_v | \bar{h}_v \delta(\ell^+ - in \cdot D) h_v | B_v \rangle, \]
where \( h_v \) is the heavy quark field. The subleading shape functions are
\[ \langle B_v | O_0(\ell^+) | B_v \rangle = f^{(2)}_0(\ell^+), \]
\[ \langle B_v | O_{\gamma}(\ell^+) | B_v \rangle = \left( \epsilon^\gamma - \frac{n^\gamma}{n \cdot v} \right) f^{(2)}_1(\ell^+), \]
\[ \langle B_v | P_{2\lambda}(\ell^+) | B_v \rangle = \epsilon^{\lambda}_\perp f^{(2)}_2(\ell^+), \]
\[ \langle B_v | O_3^{\alpha\beta}(\ell^+) | B_v \rangle = g^{\alpha\beta}_\perp f^{(4)}_3(\ell^+_1, \ell^+_2), \]
\[ \langle B_v | P_{4\lambda}^{\alpha\beta}(\ell^+_1, \ell^+_2) | B_v \rangle = -\epsilon^{\alpha\beta}_\perp \left( v_\lambda - \frac{n_\lambda}{n \cdot v} \right) f^{(4)}_4(\ell^+_1, \ell^+_2), \]
\[ \langle B_v | O_5^{\alpha\beta}(\ell^+_1, \ell^+_2, \ell^+_3) | B_v \rangle = f^{(6)}_5(\ell^+_1, \ell^+_2, \ell^+_3), \]
where \( g^{\mu\nu}_{\perp} = g^{\mu\nu} - (1/2)(n^\mu \bar{n}^\nu + n^\nu \bar{n}^\mu) \) and \( \epsilon^{\mu\nu}_\perp = (1/2)\epsilon^{\mu\nu\alpha\beta}\bar{n}_\alpha n_\beta \). The ultrasoft operators are
\[ O^{(2)}_0(\ell^+) = \int_0^{dx} e^{-\frac{1}{2}x^2-\ell^+} \int d^4y T \left[ \bar{h}_v (\tilde{x}) Y(\tilde{x}, 0) h_v (0) iO_h (y) \right], \]
\[ O^{(1)}_1(\ell^+) = \frac{1}{4} i \bar{h}_v \left[ i D^\alpha_{\mu\nu} \delta(\ell^+ - in \cdot D) \right] h_v, \]
\[ P^{(2)}_{2\lambda}(\ell^+) = \frac{i}{2} \bar{h}_v \left[ i D^\lambda_{\mu\nu} \delta(\ell^+ - in \cdot D) \right] \gamma^\lambda \gamma_5 h_v, \]
\[ O^{(3)}_{3,2}(\ell^+_1, \ell^+_2) = \frac{1}{2} \bar{h}_v \delta(\ell^+_2 - in \cdot D) \left\{ i D^\alpha_{\mu\nu} \right\} \delta(\ell^+_1 - in \cdot D) h_v, \]
\[ P^{(4)}_{4\lambda}(\ell^+_1, \ell^+_2) = -\frac{1}{2} \bar{h}_v \delta(\ell^+_2 - in \cdot D) gG^{\alpha\beta}_{\mu\nu} \delta(\ell^+_1 - in \cdot D) \gamma^\lambda \gamma_5 h_v, \]
\[ O^{(5)}_{5,6}(\ell^+_1, \ell^+_2, \ell^+_3) = \frac{1}{2} \left\{ \bar{h}_v \delta(\ell^+_3 - in \cdot D) \gamma^\alpha P_L T^A q_{\tilde{\mu}} \right\} \delta(\ell^+_2 - in \cdot \partial) \left\{ q_{\tilde{\mu}} \gamma^\alpha P_L \delta(\ell^+_1 - in \cdot D) T^A h_v \right\}, \]
where \( \tilde{x}^\mu = \bar{n} \cdot x n^\mu/2 \). Here, \( O_h \) is the NLO term in the HQET Lagrangian, \( Y \) is an ultrasoft Wilson line, \( ig^{\mu\nu}_{\perp} = [i D^\mu_{\mu\nu}, i D^\nu_{\mu\nu}] \) and \( q_{\tilde{\mu}} = (\bar{n} \gamma_{\tilde{\mu}}) / 4 q_{us} \). The operator \( O^{(5)}_{5,6} \), which appears in the definitions of \( f_{5,6} \), depends upon the light-quark flavour, \( u \) or \( s \).
APPENDIX B: HARD COEFFICIENTS

In this appendix, we present expressions for the hard coefficients in \( B \to X_u \ell \bar{\nu} \) and \( B \to X_s \ell^+ \ell^- \) [24, 32, 36]. At lowest order, we have

\[
W_i^{(0)} = h_i(p_X, m_b, \mu) \int_{0}^{p_X^+} dk^+ \mathcal{J}^{(0)}(p_X^- k^+, \mu) f_i^{(0)}(k^+ + \overline{\Lambda} - p_X^+, \mu). \tag{B1}
\]

For \( B \to X_u \ell \bar{\nu} \), we have

\[
h_1^u = \frac{1}{4} \left[ C_1^{(v)} \right]^2, \tag{B2}
\]

\[
h_2^u = \frac{(1 - u_H) \left[ (C_1^{(v)})^2 + C_1^{(v)} C_2^{(v)} + C_2^{(v)} C_3^{(v)} \right]}{(\overline{y}_H - u_H)} + \frac{(C_2^{(v)})^2}{4} + \frac{(1 - u_H)^2 \left[ (C_1^{(v)})^2 + 2C_1^{(v)} C_3^{(v)} \right]}{(\overline{y}_H - u_H)^2},
\]

\[
h_3^u = \frac{(C_1^{(v)})^2}{2m_B(\overline{y}_H - u_H)},
\]

where

\[
C_1^{(v)}(\hat{\omega}, 1) = 1 - \frac{\alpha_s(m_b)C_F}{4\pi} \left\{ 2\ln^2(\hat{\omega}) + 2\text{Li}_2(1 - \hat{\omega}) + \ln(\hat{\omega}) \left( \frac{3\hat{\omega} - 2}{1 - \hat{\omega}} \right) + \frac{\pi^2}{12} + 6 \right\},
\]

\[
C_2^{(v)}(\hat{\omega}, 1) = \frac{\alpha_s(m_b)C_F}{4\pi} \left\{ \frac{2}{(1 - \hat{\omega})} + \frac{2\hat{\omega} \ln(\hat{\omega})}{(1 - \hat{\omega})^2} \right\},
\]

\[
C_3^{(v)}(\hat{\omega}, 1) = \frac{\alpha_s(m_b)C_F}{4\pi} \left\{ \frac{(1 - 2\hat{\omega})\hat{\omega} \ln(\hat{\omega})}{(1 - \hat{\omega})^2} - \frac{\hat{\omega}}{1 - \hat{\omega}} \right\}. \tag{B3}
\]

Here, \( \hat{\omega} = \omega/m_b \).

For \( B \to X_s \ell^+ \ell^- \), we have

\[
h_1^{\ell\ell} = \frac{1}{2} \left( |C_9|^2 + |C_{10a}|^2 \right) + \frac{2\text{Re}[C_7 C_{9\ast}]}{(1 - \overline{y}_H)} + \frac{2|C_7|^2}{(1 - \overline{y}_H)^2}, \tag{B4}
\]

\[
h_2^{\ell\ell} = \frac{2(1 - u_H)}{(\overline{y}_H - u_H)} \left( |C_9|^2 + |C_{10a}|^2 + \text{Re}[C_{10a} C_{10b\ast}] \right) + \frac{|C_{10b}|^2}{2} - \frac{8|C_7|^2}{(1 - \overline{y}_H)(\overline{y}_H - u_H)},
\]

\[
h_3^{\ell\ell} = \frac{-4\text{Re}[C_{10a} C_{7\ast}]}{m_B(1 - \overline{y}_H)(\overline{y}_H - u_H)} - \frac{2\text{Re}[C_{10a} C_{9\ast}]}{m_B(\overline{y}_H - u_H)}.
\]

The full expressions for the coefficients \( C_{7,9,10a,10b} \) are given in Ref. [24]. When we ignore \( \mathcal{O}(\alpha_s(m_b)) \) corrections, they simplify to

\[
C_9 = C_{9\ast} = C_9^{\text{mix}}, \tag{B5}
\]

\[
C_7 = C_{7a} = \frac{m_b(\mu_0)}{m_B} C_7^{\text{NDR}}(\mu_0),
\]

\[
C_{10a} = C_{10a} = C_{10},
\]

\[
C_{10b} = 0,
\]

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where \( \mu_0 \sim m_b \) and

\[
C_{9}^{\text{mix}}(\mu_0) = C_{9}^{\text{NDR}}(\mu_0) + \frac{2}{9} (3C_3 + C_4 + 3C_5 + C_6) - \frac{1}{2} h(1, s) (4C_3 + 4C_4 + 3C_5 + C_6) + h\left(\frac{m_c}{m_b}, s\right) (3C_1 + 2 + 3C_3 + 4 + 3C_5 + C_6) - \frac{1}{2} h(0, s) (C_3 + 3C_4) + \mathcal{O}(\alpha_s(\mu_0)),
\]

(B6)

The function \( h(z, s) \) is given by

\[
h(z, s) = \frac{8}{9} \ln\left(\frac{\mu_0}{m_b}\right) - \frac{8}{9} \ln z + \frac{8}{27} + \frac{4}{9} \zeta - \frac{2}{9} (2 + \zeta) \sqrt{1 - \zeta} \\
\times \left[ \theta(1 - \zeta) \left( - i \pi + \ln \frac{1 + \sqrt{1 - \zeta}}{1 - \sqrt{1 - \zeta}} \right) + \theta(\zeta - 1) 2 \arctan \frac{1}{\sqrt{\zeta - 1}} \right],
\]

\[
h(0, s) = \frac{8}{27} + \frac{8}{9} \ln\left(\frac{\mu_0}{m_b}\right) - \frac{4}{9} \ln s + \frac{4}{9} i \pi,
\]

(B7)

with \( \zeta = 4z^2/s \) and \( s = q^2/m_b^2 \).

In the expressions above, \( C_{1-6}, C_{7,9}, C_{10} \) are the coefficients of the corresponding operators in the effective Hamiltonian for \( b \to s \ell^+ \ell^- \) (for which the NLL calculations were done in Refs. [42, 43]), while \( C_{9}^{\text{mix}} \) differs from \( C_{9}^{\text{eff}} \) of Ref. [42] by only an \( \mathcal{O}(\alpha_s) \) piece. Note that there is a complication in the perturbative power counting. Above the scale \( m_b \), one usually expands in \( \alpha_s \), with \( \alpha_s \log(m_W/m_b) = \mathcal{O}(1) \). Because of mixing with \( \mathcal{O}_{1,2}, C_9 \sim \log(m_W/m_b) \sim 1/\alpha_s \), whereas \( C_{7,10} \sim 1 \). However, numerically \( |C_9(m_b)| \sim C_{10} \). This problem is exacerbated by the fact that in the shape function region only the rate is calculable, not the amplitude. The solution is to use a “split matching” procedure, which decouples the scale dependence above and below \( \mu = m_b \) and thereby allows us to consider the coefficients as \( \mathcal{O}(1) \) numbers in the latter region [24].

At next-to-leading order, we have

\[
W_i^{(2)f} = \frac{h_{ij}^{ij}(\bar{n} \cdot p)}{2m_b} \int_0^{p_\ell^+} dk_1^+ \mathcal{J}^{(0)}(\bar{n} \cdot p k^+, \mu) f_0^{(2)}(k^+ + r^+, \mu) + \sum_{r=1}^{2} \frac{h_{ij}^{ij}(\bar{n} \cdot p)}{m_b} \int_0^{p_\ell^+} dk_1^+ \mathcal{J}^{(0)}(\bar{n} \cdot p k^+, \mu) f_r^{(2)}(k^+ + r^+, \mu) + \sum_{r=3}^{4} \frac{h_{ij}^{ij}(\bar{n} \cdot p)}{m_b} \int_0^{k_2^+} dk_2^+ \mathcal{J}^{(2)}(\bar{n} \cdot p k^+, \mu) f_r^{(4)}(k^+ + r^+, \mu) + \sum_{r=5}^{6} \frac{h_{ij}^{ij}(\bar{n} \cdot p)}{\bar{n} \cdot p} \int_0^{k_2^+} dk_2^+ \mathcal{J}^{(2)}(\bar{n} \cdot p k^+, \mu) f_r^{(6)}(k^+ + r^+, \mu) + \ldots,
\]

(B8)

where \( r^+ = \bar{\Lambda} - p_{\ell^+}^\perp \), \( j = 1, 2 \) and \( j' = 1, 2, 3 \). The ellipses denote terms that have jet functions \( \mathcal{J} \) that start at one-loop order or higher. (These terms are given in Ref. [32].)
When we keep the full dependence on $u_H$, the $h_{3}$'s are

$$h_{1u}^{1u} = \frac{1}{8}, \quad h_{2u}^{1u} = \frac{(1 - u_H)(2 - \bar{y}_H - u_H)}{2(\bar{y}_H - u_H)^2}, \quad h_{3u}^{1u} = \frac{1}{4m_B(\bar{y}_H - u_H)}, \quad (B9)$$

$$h_{1u}^{2u} = -\frac{1}{4}, \quad h_{2u}^{2u} = \frac{(1 - u_H)((4 - u_H)\bar{y}_H - \bar{y}_H^2 - 2)}{\bar{y}_H(\bar{y}_H - u_H)^2}, \quad h_{3u}^{2u} = -\frac{1}{2m_B(\bar{y}_H - u_H)},$$

$$h_{1u}^{3u} = -\frac{1}{4\bar{y}_H}, \quad h_{2u}^{3u} = -\frac{(1 - u_H)}{\bar{y}_H(\bar{y}_H - u_H)}, \quad h_{3u}^{3u} = -\frac{1}{2mb\bar{y}_H(\bar{y}_H - u_H)},$$

$$h_{1u}^{4u} = \frac{1}{4\bar{y}_H}, \quad h_{2u}^{4u} = -\frac{(1 - u_H)(2 - \bar{y}_H - u_H)}{\bar{y}_H(\bar{y}_H - u_H)^2}, \quad h_{3u}^{4u} = \frac{1}{2mb\bar{y}_H(\bar{y}_H - u_H)},$$

$$h_{1u}^{5u} = -\frac{1}{2}, \quad h_{2u}^{5u} = \frac{2(1 - u_H)(1 - \bar{y}_H)}{(\bar{y}_H - u_H)^2}, \quad h_{3u}^{5u} = -\frac{1}{m_b(\bar{y}_H - u_H)},$$

$$h_{1u}^{6u} = 0, \quad h_{2u}^{6u} = \frac{2(1 - u_H)^2}{(\bar{y}_H - u_H)^2}, \quad h_{3u}^{6u} = 0,$$

and the $h_{1\ell\ell}$ are

$$h_{1\ell}^{1\ell} = -\frac{4|C_{7a}|^2 - (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2}{4(1 - \bar{y}_H)^2}, \quad (B10)$$

$$h_{2\ell}^{1\ell} = \frac{(2 - \bar{y}_H - u_H)(4|C_{7a}|^2 - (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)(1 - u_H))}{(1 - \bar{y}_H)(\bar{y}_H - u_H)^2},$$

$$h_{3\ell}^{1\ell} = -\frac{\text{Re}[C_{10a}C_{9a}^*]}{m_B(\bar{y}_H - u_H)},$$

$$h_{1\ell}^{2\ell} = \frac{4|C_{7a}|^2 - (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2}{2(1 - \bar{y}_H)^2},$$

$$h_{2\ell}^{2\ell} = -\frac{2}{\bar{y}_H(\bar{y}_H - u_H)^2} \left[ 4|C_{7a}|^2 \frac{2 - \bar{y}_H^2 - \bar{y}_H u_H}{1 - \bar{y}_H} + 4\text{Re}[C_{7a}C_{9a}^*](2 - \bar{y}_H - u_H) \right.$$  
$$+ (|C_{10a}|^2 + |C_{9a}|^2)(2 - 4\bar{y}_H + \bar{y}_H^2 + \bar{y}_H u_H)(1 - u_H) \right],$$

$$h_{3\ell}^{2\ell} = \frac{2\text{Re}[C_{10a}C_{9a}^*]}{m_B(\bar{y}_H - u_H)},$$

$$h_{1\ell}^{3\ell} = -\frac{4|C_{7a}|^2 + 4\text{Re}[C_{7a}C_{9a}^*](1 - \bar{y}_H) + (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2}{2\bar{y}_H(1 - \bar{y}_H)^2},$$

$$h_{2\ell}^{3\ell} = \frac{4|C_{7a}|^2 - (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)(1 - u_H)}{\bar{y}_H(1 - \bar{y}_H)(\bar{y}_H - u_H)},$$

$$h_{3\ell}^{3\ell} = \frac{2\text{Re}[C_{10a}C_{7a}^*] + \text{Re}[C_{10a}C_{9a}^*](1 - \bar{y}_H)}{m_B\bar{y}_H(1 - \bar{y}_H)(\bar{y}_H - u_H)}.$$
\[ h_{1}^{\ell\ell} = \frac{4|C_{7a}|^2 + 4\text{Re}[C_{7a}C_{9a}^*](1 - \bar{y}_H) + (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2}{2 \bar{y}_H (1 - \bar{y}_H)^2}, \]
\[ h_{2}^{\ell\ell} = -\frac{2}{\bar{y}_H (\bar{y}_H - u_H)^2} \left[ 4|C_{7a}|^2 \frac{2 - \bar{y}_H - u_H}{1 - \bar{y}_H} + 8\text{Re}[C_{7a}C_{9a}^*](1 - u_H) + (|C_{10a}|^2 + |C_{9a}|^2)(2 - \bar{y}_H - u_H)(1 - u_H) \right], \]
\[ h_{3}^{\ell\ell} = -2 \frac{2\text{Re}[C_{10a}C_{7a}^*] + \text{Re}[C_{10a}C_{9a}^*](1 - \bar{y}_H)}{m_B \bar{y}_H (1 - \bar{y}_H)(\bar{y}_H - u_H)} , \]
\[ h_{4}^{\ell\ell} = \frac{4|C_{7a}|^2 + 4\text{Re}[C_{7a}C_{9a}^*](1 - \bar{y}_H) + (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2}{(1 - \bar{y}_H)^2}, \]
\[ h_{5}^{\ell\ell} = \frac{4 (4|C_{7a}|^2 + 4\text{Re}[C_{7a}C_{9a}^*](1 - \bar{y}_H) + (|C_{10a}|^2 + |C_{9a}|^2)(1 - \bar{y}_H)^2) (1 - u_H)}{(1 - \bar{y}_H)(\bar{y}_H - u_H)^2}, \]
\[ h_{6}^{\ell\ell} = 4 \frac{2\text{Re}[C_{10a}C_{7a}^*] + \text{Re}[C_{10a}C_{9a}^*](1 - \bar{y}_H)}{m_B (1 - \bar{y}_H)(\bar{y}_H - u_H)}, \]
\[ h_{1}^{00} = 0, \quad h_{2}^{00} = 4 \frac{4|C_{7a}|^2 + 4\text{Re}[C_{7a}C_{9a}^*](1 - u_H) + (|C_{10a}|^2 + |C_{9a}|^2)(1 - u_H)^2}{(\bar{y}_H - u_H)^2}, \]
\[ h_{3}^{00} = 0. \]

**APPENDIX C: WEIGHT FUNCTIONS**

\begin{align*}
(4) \quad K_{N}^{u} &= \frac{1}{N(u_H)} (\frac{2x_H + u_H + \bar{y}_H - 2}{\bar{y}_H - u_H}) (\frac{1 - \bar{y}_H(2\bar{y}_H - u_H - 1)}{1 + u_H - \bar{y}_H}) \\
N(u_H) &= \frac{1}{30} (1 - u_H)^2 (1 - 14u_H - 94u_H^2 - 14u_H^3 + u_H^4) - 2u_H^2 (1 - u_H^2) \log u_H \\
\kappa_{1}^{u} &= \frac{1}{2}, \quad \kappa_{2}^{u} = -1 \\
(5) \quad K_{N}^{u} &= \frac{6 \left[ -(7 - u_H)\bar{y}_H + 4x_H\bar{y}_H + 6\bar{y}_H^2 \right]}{(1 - u_H)^7} \\
\kappa_{1}^{u} &= -\frac{1}{10}, \quad \kappa_{2}^{u} = \frac{7 - u_H}{5(1 - u_H)} \\
(6) \quad K_{N}^{u} &= \frac{105 \left[ 14(1 - u_H)^2\bar{y}_H + 5(2 + 7u_H)x_H\bar{y}_H - 45x_H\bar{y}_H^2 \right]}{101 (1 - u_H)^8} \\
\kappa_{1}^{u} &= -\frac{2}{101}, \quad \kappa_{2}^{u} = \frac{109 - 4u_H}{101(1 - u_H)}
\end{align*}

**TABLE IV:** Some choices of $K^u(x_H, \bar{y}_H, u_H)$ for which the weighted integral Eq. (17) depends only on the shape functions $F, F_1$ and $F_2$. The coefficients $\kappa_{1}^{u}(u_H)$ and $\kappa_{2}^{u}(u_H)$ are defined in Eq. (37).
TABLE V: Some choices of $K^{\ell\ell}(x_H, y_H, u_H)$ for which $\kappa_2^{\ell\ell}(u_H) = \kappa_3^{\ell\ell}(u_H)$ in Eq. (37).  Here, $\mathcal{A} = -2\text{Re}[C_{10a} C_{7a}^*]$, $\mathcal{B} = \text{Re}[C_{10a} C_{9a}^*]$ and $y_H^* = y_{\text{min}} + y_{\text{max}}$. 

$^a$Note that Example (9) requires a harsher cut, e.g. $2\text{GeV}^2 \leq q^2 \leq 6\text{GeV}^2$ (rather than $1\text{GeV}^2 \leq q^2 \leq 6\text{GeV}^2$), so that it is not singular.
may be taken to be constant, in which case the integrals can be evaluated analytically.

Here, $K_{22}^\ell = -\frac{2}{N(u_H)} \int_{g_{\text{min}}}^{g_{\text{max}}} dg_H B(1 - \bar{y}_H)(\bar{y}_H - u_H)^3 \{A - B(1 - \bar{y}_H)\}$

$K_{23}^\ell = -\frac{2}{N(u_H)} \int_{g_{\text{min}}}^{g_{\text{max}}} dg_H (\bar{y}_H - u_H)^3 \{A - B(1 - \bar{y}_H)\}$

$K_{33}^\ell = 0$

**Table VI:** Some choices of $K_{\ell\ell}^r$ and the corresponding coefficients $\kappa_{2}^{\ell}(u_H)$ and $\kappa_{3}^{\ell}(u_H)$, which are defined in Eq. (37). Here, $A = -2 \text{Re}[C_{10a}^r C_{7a}^r]$, $B = \text{Re}[C_{10a}^r C_{9a}^r]$ and $\bar{y}_{H*} = \bar{y}_{\text{min}} + \bar{y}_{\text{max}}$. $C_{9a}$ may be taken to be constant, in which case the integrals can be evaluated analytically.
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