Short- and long-distance QCD effects in $B$-meson decays

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Abstract. Various exclusive and inclusive decays of $B$ mesons are studied, at present, with dedicated experiments at “$B$ factories”. In order to compete with the experimental accuracy, we need a reliable theoretical framework to compute strong interaction effects in a hadronic environment. I discuss how the separation of (perturbatively calculable) short-distance QCD effects from (non-perturbative) long-distance phenomena helps to obtain precise theoretical predictions.

INTRODUCTION

The standard model (SM) for electroweak and strong interactions of elementary particles allows for precise theoretical predictions that are in remarkable agreement with present experimental results from high-energy experiments. However, an explanation of the SM and its parameters within a more fundamental theoretical framework remains an unresolved puzzle. Many questions in this respect (e.g. a grand unification of forces and matter, the origin of neutrino masses and its relation to the quark sector, the origin of $CP$ violation, etc.) are related to the flavour sector of the SM, i.e. to the masses and coupling constants of the different quark and lepton species. $B$-meson decays are particularly useful to explore the parameters of the flavour sector in the SM or possible extensions of it. The $b$ quark has a relatively long lifetime, and provides many rare decay modes that have small branching ratios and are sensitive to details of the interactions at small distances. In this way one hopes to reveal indirect effects from physics beyond the SM even before the direct detection of new particles.

From the QCD point of view, $b$ quarks are interesting because they are the heaviest quarks that build pronounced hadronic bound states. The fact that the $b$-quark mass ($m_b \approx 5$ GeV) is large with respect to typical hadronic scales leads to new approximate symmetries that can be observed in the $b$-hadron spectrum and decays. Furthermore, the heavy-quark mass provides a scale at which the strong coupling constant is still small, and perturbative computations of short-distance effects are possible. The short-distance dynamics has to be separated from (non-perturbative) long-distance physics related to the intrinsic QCD scale $\Lambda_{\text{QCD}}$. Technically this is achieved by an operator product expansion and from an effective field theory approach, respectively. In this way, $B$-meson decays can also be used to test and improve our theoretical understanding of QCD in different dynamical regimes.
In the SM, transitions between different quark flavours are mediated by charged $W$ bosons. The relative strengths $V_{ij}$ for flavour transitions $q_i \to q_j$ define the unitary Cabbibo–Kobayashi–Maskawa (CKM) matrix, which can be parametrized in terms of 3 real angles and one $CP$-violating phase. On the other hand, flavour-changing neutral currents (FCNCs) can only be induced via loop diagrams (box or penguin topologies, see Fig. 1). As a consequence, FCNCs are sensitive to the properties of virtual heavy particles in the loops, e.g. the top quark in the SM, or new particles in extensions of it.

**Effective Hamiltonian**

As in the case of the muon decay, we may “integrate out” the heavy particles ($W, Z$ bosons, top quark, new physics) to arrive at an effective Hamiltonian (for a specific flavour transition), which has the schematic form

$$H_{\text{eff}} \propto G_F \sum_i C_i(\mu) \cdot \mathcal{O}_i + \text{terms suppressed by } 1/m_W^2.$$  \hspace{1cm} (1)

Here $G_F$ is the Fermi constant, and $C_i(\mu)$ are effective coupling constants (Wilson coefficients) that encode the short-distance dynamics from physics above the scale $\mu = O(m_W)$. The dynamics of the remaining five quark flavours, leptons, gluons and photons at energy scales below $\mu$ is described by a set of operators $\mathcal{O}_i$ (four-fermion operators, and operators that couple fermions to the chromomagnetic or electromagnetic field strength). For more details, see the review in [1].

**Matching**

The Wilson coefficients can be calculated by “matching” scattering amplitudes in the SM and the effective theory. This requires computation of QCD (and QED) radiative corrections in both cases; see Fig. 2. Since the strong coupling constant is small at the matching scale, the Wilson coefficients have a perturbative expansion,

$$C_i(m_W) = c_i^{(0)} + \frac{\alpha_s(m_W)}{4\pi} c_i^{(1)} + \ldots$$ \hspace{1cm} (2)
At present, the SM matching coefficients are known at order $\alpha_s^2$ (i.e. at two loops). Different models for physics at and above the electroweak scale yield different matching coefficients:

$$c_i^{(n)} = c_i^{(n)}(m_t, m_W, ...) + c_i^{(n)}(\text{new physics}) .$$

(3)

Therefore, experimental measurements of Wilson coefficients in weak decays test the SM and/or constrain new physics models.

**Resummation of large logarithms**

Radiative corrections to the matrix elements of the effective operators $\mathcal{O}_i$ in general would involve large logarithms $\ln m_b/\mu$, when calculated at a scale $\mu \sim m_W$. Since higher orders in perturbation theory also lead to higher powers of logarithms, the convergence of the perturbative series would be poor. This can be avoided by exploiting the fact that Wilson coefficients in the effective theory obey a renormalization-group equation, which allows the evolution of $C_i(m_W)$ to $C_i(m_b)$. At the low scale $m_b$, the large logarithms $(\ln m_b/m_W)^n$ are explicitly resummed in $C_i(m_b)$, and matrix elements of $\mathcal{O}_i$ only contain dynamics from energy scales below $\mu = m_b$. In order to derive the evolution equations, one has to calculate the anomalous-dimension matrix $\gamma_{ij}$ that describes the scale dependence and mixing of operators in the effective theory. Currently, the calculation of three-loop anomalous dimensions is being completed (for a recent contribution, see [2]).

**Example: $B \to X_s \gamma$.** A prominent phenomenological example, where the theoretical machinery of the effective-Hamiltonian approach is relevant, is the inclusive rare radiative decay $B \to X_s \gamma$. It provides a stringent test of the contributions to the Wilson coefficient $C_7$, which is related to the effective $b \to s \gamma$ vertex. The comparison between experiment and theory reads

$$\text{Exp.: } \text{BR}[B \to X_s \gamma] = (3.34 \pm 0.38) \times 10^{-4} \quad (4)$$

$$\text{SM: } \text{BR}[B \to X_s \gamma] = (3.70 \pm 0.30) \times 10^{-4} \quad (5)$$

The theoretical uncertainty is dominated by the renormalization-scheme dependence induced by the charm quark mass. The experimental and theoretical accuracy is already sufficient to put strong constraints on many new physics models (for a recent review on inclusive rare $B$ decays, see [6]).
QCD AND HEAVY-QUARK EXPANSION

Because quarks are confined, experiments can only probe weak interactions in a hadronic environment. Technically, this amounts to considering hadronic matrix elements of the effective Hamiltonian \( \mathcal{H} \). By construction, these matrix elements are sensitive to long-distance QCD dynamics, which is not accessible in perturbation theory. Nevertheless, some simplifications arise from the fact that the \( b \) quark mass is large compared to \( \Lambda_{\text{QCD}} \). On the one hand, the strong coupling constant is small, \( \alpha_s(m_b) \ll 1 \), which implies that the dynamics at distances of order \( 1/m_b \) is still perturbative. On the other hand \( \Lambda_{\text{QCD}}/m_b \ll 1 \) provides a small expansion parameter, and in the heavy-quark limit \( (m_b \to \infty) \) the number of independent unknown hadronic quantities may be less than in the general case.

**Heavy quark effective theory (HQET)**

The above observations can be formalized in terms of an effective theory (HQET) where – to first approximation – heavy \( b \) and \( c \) quarks are replaced by static colour/flavour sources, moving with a fixed velocity \( v^\mu \),

\[
h_v(x) = e^{im_Q v \cdot x} \frac{1 + \gamma^5}{2} Q(x) .
\]

(6)

The first few terms in the effective Lagrangian, resulting from integrating out “small” spinor components and “hard” quark and gluon modes (with virtualities of order \( m_Q^2 \)), reads

\[
\mathcal{L}_{\text{HQET}} = \bar{h}_v \left\{ i v \cdot D + \frac{(i \vec{D})^2}{2m_Q} + C_m(\mu) \frac{g_s}{4m_Q} \sigma_{\mu\nu} G^{\mu\nu} + \ldots \right\} h_v ,
\]

(7)

(where \( v \cdot \vec{D} = 0 \)). The first term in this expansion is independent of the heavy-quark mass and diagonal in the heavy-quark spin. As a consequence, two new symmetries arise in the heavy-quark limit: heavy-flavour symmetry reflects the fact that the soft interactions of the \( b \) or \( c \) quark become the same, once the center-of-mass motion of the heavy quark is subtracted. Heavy-quark spin symmetry is related to the fact that soft interactions do not change the spin of the heavy quark. This has important implications for phenomenological observables such as the heavy hadron mass spectrum, or heavy meson transition form factors, to be discussed below.\(^1\) The symmetries are broken by the subleading (kinetic and chromomagnetic) terms in the Lagrangian (7), as well as by perturbative matching coefficients for decay currents. For more details and references to the original literature, see for instance the review [8].

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\(^1\) HQET also applies to heavy-to-light decays as long as the energy transfer to the light quarks and gluons is small. In this kinematic regime, exclusive heavy-to-light processes are described by an effective low-energy Lagrangian for heavy and light mesons that combines HQET and chiral perturbation theory [7].
Example: Extracting $|V_{cb}|$ and $|V_{ub}|$. A well-known application of the heavy-quark mass expansion and HQET is the determination of the CKM elements $|V_{cb}|$ and $|V_{ub}|$ from $b \to c$ and $b \to u$ decays. At leading order in the $1/m_b$ expansion, the inclusive rates are just given by the partonic subprocess, which is calculable in perturbation theory (this has also been exploited in obtaining the theoretical prediction for $B \to X_s \gamma$ in [5]). Power corrections to the inclusive $B \to X$ rates start only at order $1/m_b^2$, and the corresponding non-perturbative parameters can be fitted to moments of experimentally measured inclusive decay spectra. As a result, one extracts [9]

$$
|V_{cb}|_{\text{incl.}} = 0.0421 \pm 0.0013, \\
|V_{ub}|_{\text{incl.}} = 0.00426 \pm 0.00013 \pm 0.00050.
$$

(8)

The theoretical accuracy is limited by the control on higher-order corrections in perturbation theory and in the heavy-quark expansion. For the extraction of $|V_{ub}|$ a subtle issue is how to separate the $b \to u\ell\nu$ spectrum from the $b \to c\ell\nu$ background.

**Soft-collinear effective theory (SCET)**

![FIGURE 3. Kinematics for the decay $B \to X\gamma$ at large recoil energy.](image)

The heavy-quark expansion can also be systematically applied to cases where the heavy quark decays into energetic light quarks and/or gluons, see Fig. 3. Examples are exclusive $B \to \pi\pi$ decays [10], or the endpoint spectrum in inclusive $B \to X_s \gamma$ decay [11]. In addition to the HQET field (6) and “soft” (i.e. low-energetic) light quark and gluon fields, the effective theory (SCET) for these cases includes “collinear” quark and gluon modes whose energy is proportional to the heavy-quark mass. In particular, collinear quarks in SCET are described in terms of “good” light-cone components $\xi$, in terms of which the Lagrangian that describes collinear quarks interacting with collinear gluons reads

$$
\mathcal{L}_{\text{coll}} = \bar{\xi} \left\{ i n_+ \cdot D + (i \mathcal{P}_\perp - m) \frac{1}{i n_+ \cdot D} (i \mathcal{P}_\perp + m) \right\} \frac{i}{2} \Gamma^\mu \xi.
$$

(9)

Here $n_\mu$ is a light-like vector, which is determined by the jet axis or the momentum of the outgoing hadron(s), and $n_\mu^\ast$ is another light-like vector with $(n_+ \cdot n_-) = 2$. Contributions to heavy-to-light decays from the “bad” spinor components of collinear quarks, as well as from interactions between soft and collinear particles are suppressed by either $1/m_b$ or by $\alpha_s$ (see e.g. [12, 13]). In this way SCET provides an elegant alternative to understand factorization theorems for various QCD processes [14]. Although the
physical motivation for separating short- and long-distance physics in HQET and SCET is similar, there are important differences to be noted:

- SCET gives no constraints for the light hadron spectrum.
- Interactions between energetic particles in the final state with soft spectators in the $B$ meson are mediated by particles with virtualities of order $\mu^2 = m_b \Lambda_{\text{QCD}}$, which corresponds to a second short-distance scale in addition to $m_b^2$.
- The effective Lagrangian and currents in SCET are non-local.

**Example: Forward–backward asymmetry zero in $B \to K^* \ell^+ \ell^-$.** The suppression of the “bad” spinor components for light quarks in SCET and heavy quarks in HQET leads to the reduction of independent form factors for heavy-to-light transitions at large recoil [15]. As a consequence, to first approximation, the position of the forward–backward asymmetry zero in the decay spectrum of the exclusive $B \to K^* \ell^+ \ell^-$ can be predicted in a model-independent way [16, 17] in terms of the SM Wilson coefficients $C_7$ and $C_9$ for $b \to s \gamma$ and $b \to s \ell^+ \ell^-$ transitions. First-order radiative corrections have been calculated in [18] and lead to the estimate $q_0^2 = 4.2 \pm 0.6$, where $q_0^2$ refers to the invariant mass of the lepton pair at which the forward–backward asymmetry vanishes.

**QCD AND HADRONIC EFFECTS**

After separating and calculating the short-distance QCD effects, we are left with matrix elements that encode the long-distance dynamics of quarks and gluons within hadrons. Since QCD perturbation theory cannot be applied in this case, we need alternative strategies to determine these hadronic parameters and obtain theoretical predictions for physical observables:

- extract hadronic parameters from one experimental observable, and insert them into the prediction for another;
- estimate them from non-perturbative calculations within lattice QCD;
- estimate them using QCD (light-cone) sum-rule techniques;
- tune phenomenological models to a subset of hadronic observables, and estimate other observables within that model.

In all these cases, the ultimate challenge is to give a reliable prediction of the systematic uncertainties. For the first option this apparently is a straightforward task, but only practicable if there are enough independent processes where one and the same hadronic quantity enters (which is the case, for instance, for parton distribution functions). In the lattice approach one has to deal with several (simultaneous) extrapolations: from finite lattice spacing to the continuum limit ($a \to 0$), from the “quenched approximation” ($n_f = 0$) to dynamical quarks, from quarks fitting on a finite lattice to realistic light- and heavy-quark masses. Here, progress is to be expected from improved lattice actions, increasing computer power and optimizing numerical algorithms, as well as from understanding the systematic effects for dynamical quark simulations better (see also C. Davies’ talk at this conference). QCD sum rules are based on parton–hadron
duality, which is applied to appropriate correlation functions, from which properties of the lowest contributing resonance can be extracted (for an introduction to QCD sum rules, see [19]). The suppression of contributions from higher resonances is achieved by a Borel transformation, and by relating the hadronic and the partonic representation of the spectral function above some threshold $s_0$. Among others, systematic uncertainties arise from varying the Borel parameter and the threshold parameter within a “stability window”, and from power corrections parametrized in terms of quark and gluon condensates. Finally, concerning model estimates, the only way to quantify systematic uncertainties often is to compare the predictions of sufficiently many different models.

In the following I give some examples of hadronic parameters relevant to $B$ physics. More details on the numerical estimates for various quantities can be found in [20].

### $f_B$ and $B_B$

Important phenomenological quantities that characterize the $B$ meson are the decay constant and the $B_q^0 - ar{B}_q^0$ mixing parameter. They are defined as

\[
\langle 0 | \bar{q} \gamma_\mu \gamma_5 b | B(p) \rangle = if_B q_\mu ; \quad \langle B_q^0 | \mathcal{O}(\Delta B = 2) | B_q^0 \rangle \rangle = \frac{8}{3} B_{B_q}(\mu) f_{B_q}^2 m_{B_q}^2 ,
\]

respectively, where $\mathcal{O}(\Delta B = 2)$ is a four-quark operator that induces $\bar{b}q \leftrightarrow \bar{q}b$. Numerical estimates for these quantities are summarized in Table 1 (where we considered the renormalization-group invariant quantity $\hat{B}_{B_q}$ for convenience). The ratio $\xi$ defined in that table plays an important role in the analysis of the CKM triangle. The asymmetric error in the lattice estimate arises from the different ways of performing the chiral extrapolation to realistic light-quark masses, with or without including “chiral logs”. We also remark that, in principle, $B$-meson and $D$-meson decay constants are related by HQET. However, at present, both the experimental accuracy in measuring $f_D$ and the theoretical understanding of $1/m_c$ corrections are insufficient.

### HQET parameters

In HQET the masses of the ground-state pseudoscalar ($P$) and vector ($V$) mesons to order $1/m_Q$ accuracy read

\[
\left\{ \begin{array}{c}
m_P \\
m_V 
\end{array} \right\} = m_Q + \bar{\Lambda}(m_q) + \frac{1}{2m_Q} \left\{ \begin{array}{c}
-\lambda_1(m_q) - 3\lambda_2(m_q) \\
-\lambda_1(m_q) + (1 - \delta_V(m_q)) \lambda_2(m_q)
\end{array} \right\} .
\]

In the heavy-quark limit, pseudoscalar and vector meson belong to the same spin-symmetry multiplet characterized by a residual mass term $\bar{\Lambda}(m_q)$ that only depends on the flavour of the light degrees of freedom. The parameters $\lambda_1$ and $\lambda_2$ are related to

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**Table 1.** Estimates for $B$-meson decay constant and $B_q^0 - \bar{B}_q^0$ mixing parameter from lattice QCD and QCD sum rules (numerical values taken from [20]).

|                | $f_{B_u,d}$ | $f_B$ | $\hat{B}_{B_q}$ | $B_B / B_{B_q}$ | $\xi = f_B \sqrt{\hat{B}_{B_q} / f_{B_q} \sqrt{\hat{B}_{B_q}}}$ |
|----------------|-------------|-------|-----------------|-----------------|------------------------------------------------------------------|
| **Lattice QCD** | 203(27) MeV | 238(31) MeV | 1.34(12) | 1.00(3) | 1.18(4)(12) |
| **QCD sum rules** | 208(27) MeV | 242(29) MeV | 1.67(23) | $\approx 1$ | |


the kinetic energy and chromomagnetic operator in (7), respectively. The parameters in (11) can be determined from combing information from spectroscopy, lattice, QCD sum rules, and experimental data on inclusive B-meson decays (see [20] for a summary of quantitative results). Notice that the definition of the quantities on the right-hand side in (11) is renormalization-scheme and -scale dependent.

Isgur–Wise form factor. In the heavy-quark limit, as a consequence of heavy-quark flavour and spin symmetry, all $B \to D^{(*)}$ form factors are given in terms of one Isgur–Wise function, which is normalized at zero recoil [30], e.g.,

$$\langle D|\bar{h}_v \Gamma h_v|B\rangle \propto \xi(v \cdot v') \frac{1}{m_Q^2} \left[ 1 + \frac{y'}{2} - \delta_1/m_Q^2 + \cdots \right], \quad \xi(1) = 1.$$

Perturbative QCD corrections to $\xi(1)$ can be calculated in HQET. Non-perturbative corrections to (14) start only at order $1/m_Q^2$. The estimate thus obtained for the $B \to D^*$ form factor relevant to the analysis of $|V_{cb}|$ is [20]

$$h_{A1}^{B \to D^*}(1) = \eta_A \left[ 1 + \delta_1/m_Q^2 + \delta_2/m_Q^3 + \cdots \right] = 0.91^{+0.03}_{-0.04}.$$  

The knowledge of the form factor enables one to extract the matrix element $|V_{cb}|$ from exclusive decay modes. In practice, one also needs an estimate of the slope $\xi'(1)$ to extrapolate the experimental data to zero recoil. In this way one obtains $|V_{cb}|_{\text{excl.}} =$

\[\text{...}\]

\[\text{...}\]

\[\text{...}\]

\[\text{...}\]

\[\text{...}\]

In general, the HQET parameters for the two different (would-be) spin multiplets in (11) and (12) are independent of each other. Additional constraints are obtained if one Taylor-expands all quantities around the chirally symmetric limit (i.e. vanishing quark condensate, $\langle \bar{q}q \rangle \to 0$), and if one assumes that higher-order terms in that expansion are suppressed [24, 25]. In this case the mass splitting between chiral multiplets is expected to be $\hat{\lambda}(m_q) - \bar{\lambda}(m_q) \propto |\langle \bar{q}q \rangle|$. The empirical values for this quantity are in fair agreement with phenomenological models based on spontaneous chiral symmetry breaking [26, 27, 28, 29]. Furthermore, neglecting contributions proportional to $\langle \bar{q}q \rangle$ or $m_q$ in terms that are already suppressed by $1/m_Q$, one expects

$$\lambda_1(m_q) \simeq \lambda_1'(m_q) \equiv \lambda_1, \quad \lambda_2(m_q) \simeq \lambda_2'(m_q) \equiv \lambda_2, \quad \delta_{\lambda}(m_q) \simeq \delta_{\lambda}(m_q) \equiv \delta.$$

In other words, to first approximation $1/m_Q$ corrections should not depend on the flavour and parity of the light degrees of freedom, which appears to be in good agreement with experimental data. Nevertheless, one has to keep in mind that neither $1/m_c$ nor $\langle \bar{q}q \rangle$ are very small on hadronic scales, and higher-order terms in (11)-(13) may be non-negligible.
0.0402 ± 0.0020 \[9\], which is in good agreement with the number obtained from inclusive modes \[8\].

Heavy-to-light decays. In order to extract information on heavy-to-light decays from exclusive channels, one needs (among others) the corresponding transition form factors, for instance to obtain \( C_7 \) from \( B \to K^* \gamma \) or \(|V_{ub}| \) from \( B \to \pi(\rho)\ell\nu \). In contrast to the heavy-to-heavy case, the values of heavy-to-light form factors themselves are not restricted by a simple symmetry principle. Often, one uses a simple parametrization, e.g. for the relevant form factor for \( B \to \pi\ell\nu \) decays:

\[
f_{+}^{B\rightarrow\pi}(q^2) = \frac{f_{+}^{B\rightarrow\pi}(0)}{(1-q^2/m_{B^*}^2)(1-\alpha q^2/m_{B^*}^2)},
\]

where the first term in the denominator reflects the fact that for \( q^2 \approx m_b^2 \) the form factor is dominated by the nearest vector-meson pole, and \( f_{+}^{B\rightarrow\pi}(0) = 0.2–0.3 \) and \( \alpha = 0.3–0.6 \) are fitted to lattice, QCD sum rules, or model estimates \[20\]. With increasing experimental data on the exclusive decay spectra, which would also help to test and refine different model estimates, significant improvement is expected for the determination of \(|V_{ub}| \) from exclusive decays. Currently, the result for \(|V_{ub}| \) from exclusive decays is somewhat smaller than the result from the inclusive analysis (but compatible within the rather large uncertainties) \[9\].

Other exclusive \( B \) decays into light mesons require additional non-perturbative input, as soon as radiative QCD corrections are considered. The first systematic treatment of this issue has been discussed in \[10\] for the case of non-leptonic \( B \to \pi\pi \) decays. It has been shown that (to leading order in the \( 1/m_b \) expansion and including first-order \( \alpha_s \) corrections) the decay amplitude can be factorized into perturbative short-distance coefficient functions and non-perturbative matrix elements that define heavy-to-light transition form factors, or light-cone distribution amplitudes for \( B \) mesons and pions. Comparison with experimental data from \( B \to \pi\pi \) and \( B \to \rho\ell\nu \) decays gives additional constraints on the CKM triangle (with some uncertainties coming from the parametrization of the \( B \)-meson distribution amplitude and from estimating \( 1/m_b \) corrections).

SUMMARY

The different energy scales relevant to \( B \)-meson decays imply that different dynamical aspects of QCD are probed: at the electroweak scale amplitudes are sensitive to the parameters that describe flavour transitions in the Standard Model or its possible extensions. Perturbative QCD corrections are included by matching the full theory onto an effective Hamiltonian and evolving the corresponding Wilson coefficients to scales of the order of the heavy-quark mass. At this scale, strong interactions are still perturbative. Furthermore, hadronic matrix elements can be expanded in inverse powers of the heavy-quark mass. This is formally described in terms of heavy quark effective theory (for \( b \) decays into charm quarks or low-energetic light quarks and gluons), of soft-collinear effective theory (for decays into energetic light quarks and gluons). Finally, the non-perturbative dynamics related to the intrinsic QCD scale is described in terms of
hadronic parameters that have to be extracted from experiment or estimated from lattice, QCD sum rules or models. Reducing the theoretical uncertainties from each of the above dynamical regimes, and comparing the results with experimental data from B factories, enables us to test the flavour sector of the Standard Model, to find or constrain indirect new physics contributions, and to improve our understanding of perturbative and non-perturbative QCD.

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