MPS scheme can be dramatically improved by switching to a scheme that accounts for the dominant power law dependence on the factorization scale in the operator product expansion. We introduce the “MSR scheme” which achieves this in a Lorentz and gauge invariant way. The MSR scheme has a very simple relation to MS, and can be easily used to reanalyze MS results. Results in MSR depend on a cutoff parameter $R$, in addition to the $\mu$ of MS. $R$ variations can be used to independently estimate i) the size of power corrections, and ii) higher order perturbative corrections (much like $\mu$ in MS). We give two examples at three-loop order, the ratio of mass splittings in the $B^*\to B$ and $D^*\to D$ systems, and the Ellis-Jaffe sum rule as a function of momentum transfer $Q$ in deep inelastic scattering. Comparing to data, the perturbative MSR results work well even for $Q \sim 1$ GeV, and power corrections are reduced compared to MS.

Introduction and Formalism

The operator product expansion (OPE) is an important tool for QCD. In hard scattering processes two important scales are $Q$, a large momentum transfer or mass, and $\Lambda_{\text{QCD}}$, the scale of nonperturbative matrix elements. The Wilsonian OPE introduces a factorization scale $\Lambda'$, where $\Lambda_{\text{QCD}} < \Lambda' < Q$, and expands in $\Lambda_{\text{QCD}}/Q$. Consider a dimensionless observable $\sigma$ whose OPE is

$$
\sigma = C_0^{\text{W}}(Q, \Lambda', \%0) + C_1^{\text{W}}(Q, \Lambda', \%1) + \ldots
$$

The $C_0^{\text{W}}$ are dimensionless Wilson coefficients containing contributions from momenta $k > \Lambda'$ with perturbative expansions in $\alpha_s$, and $\theta_0^{\text{W}}(\%0) = (\%0)_{\text{W}}$ are non-perturbative matrix element with mass dimensions 0 and $p$, containing contributions from $k < \Lambda'$. If $C_{n}^{\text{W}}(Q, \Lambda')$ are expanded they contain an infinite series of terms, $(\Lambda'/Q)^n$, modulo $\ln^m(\Lambda'/Q)$ terms, and this reflects the fact that $C_{0,1}^{\text{W}}$ only include contributions from momenta $k > \Lambda'$. The Wilsonian OPE provides a clean separation of momentum scales, but can be technically challenging to implement. In particular, it is difficult to define $\Lambda'$ and retain gauge symmetry and Lorentz invariance, and perturbative computations beyond one-loop are atrocious.

A popular alternative is the OPE with dimensional regularization and the $\text{MS}$ scheme, which preserves the symmetries of QCD and provides powerful techniques for multiloop computations. In this case Eq. (1) becomes

$$
\sigma = C_0(Q, \mu)\theta_0(\mu) + C_1(Q, \mu)\theta_1(\mu) + \ldots
$$

where $\mu$ is the renormalization scale and bars are used for $\text{MS}$ quantities. In $\text{MS}$ the $C_i$ are simple series in $\alpha_s$,

$$
C_i(Q, \mu) = 1 + \sum_{n=1}^{\infty} b_n\left(\frac{\mu}{Q}\right)^n \left[\frac{\alpha_s(\mu)}{4\pi}\right]^n,
$$

with coefficients $b_n(\mu/Q) = \sum_{k=0}^{n} b_{nk} \ln^k(\mu/Q)$ containing only $\ln \mu/Q$. We will always rescale $\sigma$ and the matrix elements $\theta_i$ such that $C_i = 1$ at tree level. In $\text{MS}$ all power law dependence on $Q$ is manifest and unique in each term of Eq. (3). Also simple renormalization group equations in $\mu$, like $d\ln C_0(Q, \mu)/d\ln \mu = [\gamma_s(\mu)]$, can be used to sum large logs in Eq. (3) if $Q \gg \Lambda_{\text{QCD}}$.

$C_i(Q, \%0)$ and $C_j(Q, \%0)$ are related to each other in perturbation theory, so Eqs. (1) and (3) are just the same OPE in two different schemes. The renormalization scale $\mu$ in $\text{MS}$ plays the role of $\%0$. This is precisely true for logarithmic contributions, $\ln \mu \leftrightarrow \ln(\%0)$, and here the Wilsonian picture of scale separation in $C_j$ and $\theta_i$ carries over. However, the same is not true for power law dependences on $\%0$. $\text{MS}$ integrations are carried out over all momenta, so the $C_i$ actually contain some contributions from arbitrary small momenta, and the $\theta_i$ have contributions from arbitrary large momenta. For the power law terms there is no explicit scale separation in $\text{MS}$, and correspondingly no powers of $\mu$ appear in Eq. (3). While this simplifies higher order computations, it is known to lead to factorial growth in the perturbative coefficients. For $C_0$, one has $b_{n+1}(\mu/Q) \approx (\mu/Q)^n n! [2\beta_0/p]^n Z$ at large $n \%0$, for constant $Z$. In practice this sometimes leads to poor convergence already at one or two loop order in QCD. This poor behavior is canceled by corresponding instabilities in $\theta_1$, and is referred to as an order-$p$ infrared renormalon in $C_0$ canceling against an ultraviolet renormalon in $\theta_1$. The cancellation reflects the fact that the $\text{MS}$ OPE does not strictly separate momentum scales.

The OPE can be converted to a scheme that removes this poor behavior, but still retains the simple computational features of $\text{MS}$. Consider defining a new “R-scheme” for $C_0$ by subtracting a perturbative series

$$
C_0(Q, R, \mu) = C_0(Q, \mu) - \delta C_0(Q, R, \mu),
$$

$$
\delta C_0(Q, R, \mu) = \left(\frac{R}{Q}\right)^{\beta} \sum_{n=1}^{\infty} d_n\left(\frac{\mu}{R}\right)^n \left[\frac{\alpha_s(\mu)}{4\pi}\right]^n,
$$

with $d_n(\mu/R) = \sum_{k=0}^{n} d_{nk} \ln^k(\mu/R)$. If for large $n$ the coefficients $d_n$ are chosen to have the same behavior as $b_n$,
so \( d_{n+1}(\mu/Q) \simeq (\mu/R)^n n!|2\beta_0/p|^n Z \), then the factorial growth in \( C_0(Q, \mu) \) and \( \delta C_0(Q, R, \mu) \), cancel,

\[
C_0(Q, R, \mu) \sim \left[ \frac{\mu^p}{Q^p} - \frac{R^p}{Q^p} \right] \sum_n n! \left[ \frac{2\beta_0}{p} \right]^n Z. \quad (5)
\]

Thus the R-scheme introduces power law dependence on the cutoff, \((R/Q)^p\), in \( C_0(Q, R, \mu) \), which captures the dominant \((N'/Q)^p\) behavior of the Wilsonian \( C_0^W \). In practice this improves the convergence in \( C_0 \) even at low orders in the \( \alpha_s \) series. The dominant effect of this change is compensated by a scheme change to \( \theta_1 \), \( \theta_1(\mu) = \theta_1(R, \mu) - |Q^p| Q \hat{C}_0(Q, R, \mu) \delta \theta_0(\mu) \), and the new \( \theta_1 \) will exhibit improved stability. In the R-scheme the OPE becomes

\[
\sigma = C_0(Q, R, \mu) \theta_0(\mu) + C_1(Q, R, \mu) \theta_1(R, \mu)/Q^p + C_1'(Q, R, \mu) \theta_1'(R, \mu)/Q^p + \ldots, \quad (6)
\]

where \( \theta_1' = |Q^p| Q \hat{C}_0(\mu) \delta \theta_0(\mu) \) and \( C_1' = 1 - C_1 \sim \alpha_s \). Both \( C_0 \) and \( \theta_1 \) are free of order-\( p \) renormalons. The series in \( C_1' \) \( \theta_1' \) is Borel summable. In all examples below \( \theta_0 \) is also renormalon free. The above procedure may be repeated for higher renormalons and the higher power terms in the OPE indicated by ellipses, to improve the behavior of these terms as well. At the order at which we work, we will consistently set \( C_1 = 1 \) and drop \( \theta_1' \) in the following.

To set up an appropriate R-scheme it remains to define the \( d_n \). In the renormalon literature such scheme changes are well known for masses. For OPE predictions a “renormalon subtraction” (RS) scheme has been implemented in Ref. [1]. In RS an approximate result for the residue of the leading Borel renormalon pole is used to define the \( d_n \), which adds a source of uncertainty.

For our analysis we define the “MSR” scheme for \( C_0 \) by simply taking the coefficients of the subtraction to be exactly the \( \overline{\text{MS}} \) coefficients. In general it is more convenient to use \( \ln C_0 \) rather than \( C_0 \), since this simplifies renormalization group equations. Writing the series as

\[
\ln C_0(Q, R, \mu) = \sum_{n=1}^{\infty} a_n(\mu/Q) \left[ \frac{\alpha_s(\mu)}{(4\pi)} \right]^n, \quad (7)
\]

with \( a_n(\mu/Q) = \sum_{k=0} a_n k^k \mu/Q \) we define the MSR scheme by the series

\[
\ln C_0(Q, R, \mu) = \sum_{n=1}^{\infty} \left\{ a_n \left( \frac{\mu}{Q} \right) - \frac{R^p}{Q^p} a_n \left( \frac{\mu}{R} \right) \right\} \left[ \frac{\alpha_s(\mu)}{(4\pi)} \right]^n. \quad (8)
\]

This definition still cancels the order-\( p \) renormalon for large \( n \), as in Eq. [8]. It yields the very simple relation

\[
C_0(Q, R, \mu) = C_0(Q, \mu) \left[ \frac{C_0(R, \mu)}{C_0(Q, \mu)} \right]^{-\left( R/Q \right)^p}, \quad (9)
\]

which must be expanded order-by-order in \( \alpha_s(\mu) \) to remove the renormalon. Thus the coefficient \( C_0(Q, R, \mu) \) for the MSR scheme is obtained directly from the \( \overline{\text{MS}} \) result. Note \( C_0(Q, Q, \mu) = 1 \) to all orders. The appropriate \( p \) is obtained from the \( \overline{\text{MS}} \) OPE by \( p = \text{dimension}(\theta_0) \).

The appropriate values for \( R \) in Eqs. [4] are constrained by power counting and the structure of large logs in the OPE. The power counting \( \theta_1 \sim \Lambda_{\text{QCD}} \) implies \( \theta_1 \sim \Lambda_{\text{QCD}} \), so for the matrix element we need \( R \sim \mu \gtrsim \Lambda_{\text{QCD}} \) (meaning a larger value where perturbation theory for the OPE still converges), which minimizes \( \ln(\mu/\Lambda_{\text{QCD}}) \) and \( \ln(\mu/R) \) terms in \( \theta_1(R, \mu, \Lambda_{\text{QCD}}) \). On the other hand, \( C_0(Q, R, \mu) \) has \( \ln(\mu/R) \) and \( \ln(\mu/R) \) terms, and for \( R \sim \Lambda_{\text{QCD}} \) no choice of \( \mu \) avoids large logs. For \( R \sim \mu \sim Q \) we can minimize the logs in \( C_0(Q, R, \mu) \), but not in \( \theta_1(R, \mu, \Lambda_{\text{QCD}}) \). When the OPE is carried out in \( \overline{\text{MS}} \) this problem is dealt with using a \( \mu \)-RGE to sum large logs between \( Q \) and \( \Lambda_{\text{QCD}} \). For MSR we must use \( R \)-evolution, an RGE in the \( \overline{\text{MS}} \) variable \( \overline{\text{MS}} \). The appropriate R-RGE is formulated with \( \mu = R \) to ensure there are no logs in the anomalous dimension. For \( C_0 \),

\[
\frac{R d}{dR} \ln C_0(Q, R, R) = \gamma[\alpha_s(R)] - \left( \frac{R}{Q} \right)^p \gamma[\alpha_s(R)]. \quad (10)
\]

where

\[
\gamma[\alpha_s] = \sum_{n=0}^{\infty} \gamma_n[\alpha_s(R)/4\pi]^{n+1}, \quad (11)
\]

and \( \gamma_0 = -\alpha_0^2/\beta_0 / (2\pi \alpha_s) \). Here \( \gamma_\alpha = \sum_{n=0} \gamma_n[\alpha_s(R)/4\pi]^{n+1} \) are the \( \overline{\text{MS}} \) and \( \overline{\text{MS}} \) anomalous dimensions. The choice in Eq. [8] keeps Eq. [10] simple. In cases where \( \gamma \) is absent we expect Eq. [10] to converge to lower scales due to the \((R/Q)^p\) factor multiplying \( \gamma \). For \( R_1 > R_0 \) the solution of Eq. [10] is \([U_\alpha = U_\alpha(R_1, R_0)] \)

\[
C_0(Q, R_0, R_0) = C_0(Q, R_1, R_1) U_R(R_1, R_0) U_\mu(R_1, R_0) \mu_\mu, \quad (12)
\]

where \( U_\mu \) is a usual \( \overline{\text{MS}} \) evolution factor and \( U_R \) is the \( R \)-evolution. For \( p = 1 \) the complete solution for \( U_R \) was obtained in Ref. [1]. It is straightforward to generalize this to any \( p \). At \( N^{k+1} \mathrm{LL} \) order the (real) result is

\[
U_R(Q, R_1, R_0) = \exp \left\{ \left( \frac{\Lambda_{\text{QCD}}(Q)}{Q} \right)^{k} \sum_{j=0}^k S_j \left[ \hat{p}_j \mu \right]^\mu \right\}, \quad (13)
\]

with \( \Gamma(c, t) \) the incomplete gamma function and \( t_{0,1} = -2\pi/(\beta_0 \alpha_s(R_0, 1)) \). \( \Lambda_{\text{QCD}}^{(0)} = Re^t, \Lambda_{\text{QCD}}^{(1)} = Re^t(\beta_1) \), and \( \Lambda_{\text{QCD}}^{(2)} = Re^t(\beta_1^2 - \beta_2^2) \) are evaluated at a large reference \( R \) with \( t = -2\pi/(\beta_0 \alpha_s(R_1)) \), and \( b_1 = \beta_1/(2\beta_0^2), b_2 = (\beta_2^2 - \beta_3^2)/(4\beta_0^2), b_3 = (\beta_2^2 - 2\beta_0 \beta_3)/(8\beta_0^2) \). Defining \( \gamma_0 = \gamma_0[\alpha_s(Q)/2\pi]^{n+1} \), the coefficients of \( U_R \) needed for the first three orders of \( R \)-evolution are

\[
S_0 = \gamma_0, \quad S_1 = \gamma_0 - (\hat{b}_1 + \hat{b}_2) \gamma_0, \quad S_2 = \gamma_0 - (\hat{b}_1 + \hat{b}_2) \gamma_0 + \left[ (1 + \hat{b}_1) b_2 + \hat{b}_2(\hat{b}_2 + \hat{b}_3) \right] \gamma_0. \quad (14)
\]

Eq. [10] becomes \( C_0(Q, R_1, R_1) U_R(Q, R_1, R_0) U_\mu(R_1, R_0) / \theta_0(R_0) + \theta_1(R_0, R_0)/Q^p \), and this result sums logs between \( R_1 \sim Q \) and \( R_0 \sim \Lambda_{\text{QCD}} \). This gives natural \( R \) scales for coefficients and matrix elements in the OPE.
Heavy Meson Mass Splittings in MSR

The MS OPE for the mass-splitting of heavy mesons, 
\[ \Delta m_H^2 = n_H^2 - m_H^2 \] for \( H = B, D, \), is
\[ \Delta m_H^2 = C_G(m_Q, \mu) \mu^2_G(\mu) + \sum_i C_i(m_Q, \mu) 2\rho_i(\mu)/(3m_Q) + \mathcal{O}(\Lambda^3_{QCD}/m_Q^2), \]
where \( m_Q = m_b \text{ or } m_c. \) Here \( \mu^2_G = -(B_i h_i g^{\nu\rho} G^{\mu
u} h_i B_c)/3 \) is the matrix element of the chromomagnetic operator, and \( \rho_i^3 \) for \( i = \pi G, A, LS, A_G \) are \( \mathcal{O}(\Lambda^3_{QCD}) \) matrix elements
\[ \rho_i^3 = (3/2)\Lambda \mu^2_G(\mu). \]
The order of our analysis tree level values for the \( C_i \) soffice, with 
\[ \Sigma_\rho(\mu) = (2/3)[\rho^3_G(\mu) + \rho^3_\Lambda(\mu) - \rho^3_L(\mu) + \rho^3_A(\mu)], \]
\[ \Delta m_H^2 = C_G(m_Q, \mu) \mu^2_G(\mu) + \Sigma_\rho(\mu)/m_Q + \ldots. \] (14)

Taking the ratio of mass splittings \( r = \Delta m_B^2/\Delta m_D^2 \) gives
\[ r = \frac{C_G(m_b, \mu)}{C_G(m_c, \mu)} + \frac{\Sigma_\rho(\mu)}{\mu^2_G(\mu)} \left( \frac{1}{m_b} - \frac{1}{m_c} \right) + \ldots. \] (15)

The first term in this OPE gives a purely perturbative prediction for \( r. \) The MS is known to suffer from a \( \mathcal{O}(\Lambda^3_{QCD}/m_Q) \) infrared renormalon ambiguity \[ \] with a corresponding ambiguity in \( \Sigma_\rho(\mu). \) The three-loop computation of Ref. \[ \] yields, \( r = 1 - 0.1113 \alpha_s - 0.0780 \mu_c - 0.0755 \mu_b \) at fixed order with \( m_c, \mu = m_b, \) and \( r = (0.851)^{\text{LL}} + (-0.0696)^{\text{NNLL}} + (-0.0980)^{\text{NNNLL}} \) in RGE-improved perturbation theory, with no sign of convergence in either case. In MS these leading predictions are unstable due to the \( p = 1 \) renormalon in \( C_G. \)

Let's examine the analogous result in the MSR scheme

\[ \Delta m_H^2 = C_G(m_Q, R, \mu) \mu^2_G(\mu) + \Sigma_\rho(\mu)/m_Q + \ldots. \] (16)

Since \( p = 1 \) the MSR definition in Eq. \[ \] gives
\[ C_G(m_Q, R, \mu) = \bar{C}_G(m_Q, \mu)[\bar{C}_G(\mu, R, \mu)]^{-1}/R/m_Q, \] (17)
where \( \bar{C}_G(\mu) \) is obtained from Ref. \[ \] and we expand in \( \alpha_s(\mu). \) The OPE in MSR at a scale \( R \geq \Lambda_{QCD} \)
\[ r = \frac{C_G(m_b, R_0, R_0)}{C_G(m_c, R_0, R_0)} + \frac{\Sigma_\rho(R_0, R_0)}{\mu^2_G(R_0)} \left( \frac{1}{m_b} - \frac{1}{m_c} \right). \] (18)

Large logs in \( C_G(m_Q, R_0, R_0) \) can be summed in the R-RGE in Eqs. \[ \] for simplicity we integrate out the \( b \) and \( c \)-quarks simultaneously at a scale \( R_1 \approx \sqrt{m_b m_c} \gg R_0 \approx \Lambda_{QCD}. \) This scale for \( R_1 \) keeps \( \ln(R_1/m_b, m_c) \) small. With R-evolution and \( U_R \) from Eq. \[ \] we have
\[ r = \frac{C_G(m_b, R_0, R_0)}{C_G(m_c, R_0, R_0)} U_R(m_b, m_c, R_0) + \frac{\Sigma_\rho(R_0, R_0)}{\mu^2_G(R_0)} \left( \frac{1}{m_b} - \frac{1}{m_c} \right). \] (19)

This expression is independent of \( R_1 \) and \( R_0. \) Order-by-order, varying \( R_1 \) about \( \sqrt{m_b m_c} \) yields an estimate of higher order perturbative uncertainties, much like varying \( \mu \) in MS. For \( R_0 \) the dependence cancels between the
\[ r = 0.860 \pm (0.005) \Sigma_\rho \pm (0.008)_{\text{pert}} \ldots \] (20)

This estimate for the \( \Sigma_\rho \) power correction in MSR is in good agreement with experiment, \( r_{\text{exp}} = 0.886 \) \((D_{u,d}^{(s)}, B_{u,d}^{(s)}) \) and \( 0.854 \) \((D_{s}^{(s)}, B_{s}^{(s)}) \). MSR achieves a convergent perturbative prediction for \( r \) at leading order in the OPE, and a \( 1/m_Q \) power correction of moderate size, \( \sim 0.065, \) significantly smaller than the dimensional analysis estimate of \( \Lambda_{QCD}/(m_b - 1/m_b) \sim 0.2 \) in MS.

Ellis-Jaffe sum rule in MSR

In MS the Ellis-Jaffe sum rule \[ \] for the proton in DIS with momentum transfer \( Q \) is
\[ M_1(Q) = [C_B(Q, \mu) \theta_B + C_0(Q, \mu) a_0/9] + \theta_1(Q)/Q^2. \] \( C_{B_0} \) are
FIG. 2: Perturbative results for the Ellis-Jaffe sum rule in the MSR, RS, and ∆S schemes, at leading order in 1/Q. For all curves the one parameter, ̂a₀, is fixed by data at Q ≃5 GeV.

FIG. 3: Uncertainty estimates in the MSR scheme and ∆S scheme for the Ellis-Jaffe sum rule at leading order in 1/Q.

known at 3 loops [10]. The two leading order terms are written so that both coefficients and matrix elements are separately ̂µ-independent: θ₀ = g_A/12 + a_s/36 is given by the axial couplings g_A = 1.2694 and a_s = 0.572 for the nucleon and hyperon, while ̂a₀ is a ∆S independent ∆S matrix element. θ₁ denotes all 1/Q² power corrections with their Wilson coefficients at tree level. The ∆S coefficients are affected by a p = 2 renormalons [11], which is removed in the MSR scheme. Eq. (10) gives [i = B, 0]

\[ C_i(Q, R, R) \equiv \tilde{C}_i(Q, R)[\tilde{C}_i(R, R)]^{-R^2/Q^2}. \]  

With R-evolution the MSR OPE prediction is

\[ M_1(Q) = [C_B(Q, R_1, R_1)U_R^B(Q, R_1, R_0)\theta_B + C_0(Q, R_1, R_1)U_R^0(Q, R_1, R_0)\hat{a}_0/9] + \theta_1(R_0, R_0)/Q^2, \]  

where U_R^B,0 are given by Eq. (12) with p = 2 and the corresponding (a_0)B,0 determine the appropriate (γ_n)^B,0.

Figures 2 and 3 show perturbative predictions for the Ellis-Jaffe sum rule at leading power in 1/Q, compared with proton data from Ref. [12]. We use a_s(4 GeV) = 0.2282, and the 4-loop β with 4 flavors. In Fig. 2, we show order-by-order results for the ∆S scheme at μ = Q, and for the resummed MSR scheme with R₁ = Q and R₀ = 0.9 GeV. We fix ̂a₀ = 0.141 so that ∆S and MSR agree with the data at Q ≃5 GeV. ∆S agrees well with the data for large Q, but turns away at Q ≲2 GeV and no longer converges. In contrast the MSR results still converge quickly and exhibit excellent agreement with the data over a wide range of Q values. The NLL MSR result already has the right curvature, and the NNLL and N^{3}LL curves further improve the agreement. We also include predictions in the RS scheme with subtraction scale η = 1.0 GeV from Fig. 3 of Ref. [13], which improve slightly over the ∆S results, but may not be capturing the dominant power law dependence on the factorization scale. In Fig. 3, we show uncertainties for three loop results in the ∆S and MSR schemes. The dashed red curve is the MSR prediction, and the blue band estimates the higher-order perturbative uncertainties varying μ in the range μ^{min}(Q) < μ < 2Q. For Q > 1.5 GeV, μ^{min} = Q/2, while for Q < 1.5 GeV, μ^{min} = 1.3Q/(1.1 + Q/(1 GeV)). The red solid line is the MS prediction, the red band is the perturbative uncertainty from varying R₁ in the same range as was done for μ in ∆S, and the green band estimates the 1/Q² power correction by varying R₀ = 0.7 to 1.2 GeV. Fig. 3 implies −0.01 GeV² ≲ θ₁(R₀, R₀) ≲ 0.01 GeV² in MSR, which is a much smaller power correction than the ∼0.1 GeV² estimate obtained from naive dimensional analysis in ∆S.

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