

On the ultimate uncertainty of the top quark pole mass

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Abstract

We combine the known asymptotic behaviour of the QCD perturbation series expansion, which relates the pole mass of a heavy quark to the $\overline{\text{MS}}$ mass, with the exact series coefficients up to the four-loop order to determine the ultimate uncertainty of the top-quark pole mass due to the renormalon divergence. We perform extensive tests of our procedure by varying the number of colours and flavours, as well as the scale of the strong coupling and the $\overline{\text{MS}}$ mass, and conclude that this uncertainty is around 70 MeV. We further estimate the additional contribution to the mass relation from the five-loop correction and beyond to be 250 MeV.

1 Introduction

The top quark mass is a fundamental parameter of the Standard Model (SM). Due to its large size, it has non-negligible impact in the precision tests of the SM. After the discovery of the Higgs boson and the measurement of its mass, the values of the W and top mass are strongly correlated, such that a precise determination of both parameters would lead to a SM test of unprecedented precision [1]. Indeed, there is presently some tension between the value of the top mass 177 ± 2.1 GeV fitted from electroweak data and from its direct measurement [1], for which the combination of the Tevatron and LHC data yields the 1.6σ lower value of $173.34 \pm 0.27 \pm 0.71$ GeV [2]. The value of the top mass is also crucial to the issue of stability of the SM vacuum [3]. Higgs quartic coupling decreases at high scales, eventually becoming negative. This evolution is very sensitive to the top mass value. For example, a top mass near 171 GeV would imply that the quartic coupling may vanish at the Planck scale, rather than turn negative.

The standard direct determination of the top mass at hadron colliders, being based upon observables that are related to the mass of the system comprising the top decay products, are quoted as measurements of the pole mass. On the other hand, it seems more natural to use the $\overline{\text{MS}}$ mass in both precision electroweak observables and in vacuum stability studies. In [4] the relation between the $\overline{\text{MS}}$ and pole mass for a heavy quark (the “mass conversion formula” from now on) has been computed to the fourth order in the strong coupling α_s . Assuming the value of 163.643 GeV for the top-quark $\overline{\text{MS}}$ mass $\overline{m}_t = m_t(\overline{m}_t)$, and assuming $\alpha_s^{(6)}(m_t) = 0.1088$, we have [4]

$$m_P = 163.643 + 7.557 + 1.617 + 0.501 + (0.195 \pm 0.005) \text{ GeV} \quad (1.1)$$

for the series expansion of the mass conversion formula. The last term from the fourth order correction is less than one half of the third order one.

It is also known that the mass conversion formula is affected by infrared (IR) renormalons [5–7]. This means that there are factorially growing terms of infrared origin in the perturbative expansion, such that the expansion starts to diverge at some order. If the series is treated as an asymptotic expansion, the ambiguity in its resummation is of order of a typical hadronic scale. Because of this, it is often stated that the ultimate accuracy of top pole mass cannot be below a few hundred MeV. One of the goals of this work is to make this estimate more precise.

It is remarkable that the perturbative relation between the pole and $\overline{\text{MS}}$ mass of a heavy quark appears to be dominated by the leading infrared renormalon already in low orders [8, 9]. This observation was used in previous work [10, 11], and more recently in [12, 13] to estimate the unknown normalization of the leading IR renormalon, and mostly applied in the context of bottom physics. In the context of top physics, the importance of this issue was raised recently in [14]. The purpose of this work is to combine the newly available four-loop coefficient [4] in the mass conversion formula with the known structure of the first infrared renormalon singularity [7] to determine the normalization constant and discuss its impact on top physics. We also perform an analysis of the dependence on the number of colours and flavours, which is by itself of interest, and stability tests with respect

to variations of the scale of the strong coupling and $\overline{\text{MS}}$ mass. This leads to an expression for the mass conversion factor including an estimate of the contributions beyond four loops, and an estimate of the irreducible error.

2 Reminder

The renormalon divergence is a manifestation of the fact that the mass conversion formula, while infrared finite is sensitive to small loop momentum. In the case of the pole mass this sensitivity is particularly strong, namely linear, resulting in rapid divergence of the perturbative expansion, and an infrared sensitivity of order Λ_{QCD} [5, 6]. The ambiguity in defining the pole mass is therefore of similar size. This is not surprising as the pole mass of a quark is not an observable due to confinement and the difference with the physical heavy meson masses is also of order Λ_{QCD} . Unlike other heavy quarks, the top quark decays on hadronic time scales, and thus the propagator pole position acquires an imaginary part. The renormalon divergence is not altered [15] by the fact that the top quark is unstable with a width larger than Λ_{QCD} and hence does not form bound states. The finite width simplifies the perturbative treatment of top quarks, since it provides a natural IR cut-off, and there exists no quantity for which the pole mass would ever be relevant. But the infrared sensitivity of the QCD corrections to the mass conversion factor, which causes the divergence, remains unaffected by the width.

Slightly more technically, the divergence arises from logarithmic enhancements of the loop integrand. Heuristically, this can be understood by noticing that the running coupling evaluated at the scale l of the loop momentum has the expansion

$$\alpha_s(l) = \frac{1}{b_0 \ln l^2 / \Lambda_{\text{QCD}}^2} = \frac{\alpha_s(m)}{1 - \alpha_s(m) b_0 \ln m^2 / l^2} = \sum_1^{\infty} \alpha_s^n(m) b_0^n \ln^n \frac{m^2}{l^2}. \quad (2.1)$$

The IR contribution to the last loop integration in the $(n+1)$ -loop order then takes the form

$$\delta m^{(n+1)} \propto \alpha_s^{n+1}(m) \int^m dl b_0^n \ln^n \frac{m^2}{l^2} = m (2b_0)^n \alpha_s^{n+1}(m) n!. \quad (2.2)$$

With this behaviour the series of mass corrections reaches a minimal term of order

$$\begin{aligned} m (2b_0)^n \alpha_s^{n+1} n! &\approx m \alpha_s n^{-n} (\sqrt{2\pi} n^{n+1/2} e^{-n}) \approx m \sqrt{\frac{\pi \alpha_s}{b_0}} \exp\left(-\frac{1}{2b_0 \alpha_s}\right) \\ &\approx \sqrt{\frac{\pi \alpha_s}{b_0}} \Lambda_{\text{QCD}}, \end{aligned} \quad (2.3)$$

when $n \approx 1/(2b_0 \alpha_s)$ and then diverges. Asymptotic expansions can sometimes be summed using the Borel transform. Given a power series

$$f(\alpha_s) = \sum_{n=1}^{\infty} c_n \alpha_s^n, \quad (2.4)$$

the corresponding Borel transform is defined by

$$B[f](t) = \sum_{n=0}^{\infty} c_{n+1} \frac{t^n}{n!}. \quad (2.5)$$

The Borel integral

$$\int_0^\infty dt e^{-t/\alpha_s} B[f](t) \quad (2.6)$$

has the same series expansion as $f(\alpha_s)$ and provides the exact result under suitable conditions. However, for the case of (2.2), where $c_{n+1} = (2b_0)^n n!$, the Borel integral

$$\int_0^\infty dt e^{-t/\alpha_s} \frac{1}{1 - 2b_0 t} \quad (2.7)$$

cannot be performed because of the pole at $t = 1/(2b_0)$. We can introduce some prescription for handling the pole in the integral, as, for example, the principal value prescription. Whether or not this reconstructs the exact result, an ambiguity remains, quantified by the imaginary part of the integral when going above or below the singular point. A commonly used procedure is to define this ambiguity to be equal to the imaginary part of the integral divided by Pi (see, e.g., [16], section 5.2). For (2.7), this yields

$$\Lambda_{\text{QCD}}/(2b_0). \quad (2.8)$$

In the range of α_s values considered in this paper, the ambiguity is close to the size of the smallest term in (2.3).¹

It can be shown [7] that while the precise asymptotic behaviour of the mass conversion formula differs from the simple ansatz employed in this section for illustration, as discussed below, the ambiguity is exactly proportional to Λ_{QCD} , which evaluates to about 250 MeV in the $\overline{\text{MS}}$ scheme. In the remainder of this work, we aim to quantify the proportionality factor.

3 The leading pole mass renormalon

We write the perturbative expansion of the mass conversion formula as

$$m_P = m(\mu_m) \left(1 + \sum_{n=1}^{\infty} c_n(\mu, \mu_m, m(\mu_m)) \alpha_s^n(\mu) \right). \quad (3.1)$$

Here $\alpha_s(\mu)$ is the $\overline{\text{MS}}$ coupling in the n_l light flavours theory, and $m(\mu_m)$ stands for the $\overline{\text{MS}}$ mass evaluated at the scale μ_m . (In the following we will consider different scale choices for the heavy quark mass and the strong coupling constant). We also use \overline{m} to denote the $\overline{\text{MS}}$ mass evaluated self-consistently at a scale equal to the mass itself, i.e.

$$\overline{m} = m(\overline{m}). \quad (3.2)$$

¹Note, however, the different parametric dependence on α_s of (2.3) and (2.8). The correct dependence is that of (2.8), for the following reason: The typical width of the region where the minimal term is attained grows parametrically as $\sqrt{1/(2b_0\alpha_s)}$. The accuracy of an asymptotic series is better estimated by the minimal term times the factor accounting for the number of terms in this region, which makes (2.3) parametrically consistent with (2.8). Numerically, this factor turns out to be of order one for the applications considered in this paper, as will be confirmed in section 4 below. In case of doubt, the estimate from the ambiguity of the Borel integral should be the preferred choice.

The leading IR renormalon divergence implies the following large- n behaviour of the perturbative coefficients [7] (and [16]², eq. (5.90))

$$c_n(\mu, \mu_m, m(\mu_m)) \xrightarrow{n \rightarrow \infty} N c_n^{(\text{as})}(\mu, m(\mu_m)) \equiv N \frac{\mu}{m(\mu_m)} \tilde{c}_n^{(\text{as})}, \quad (3.3)$$

where

$$\tilde{c}_{n+1}^{(\text{as})} = (2b_0)^n \frac{\Gamma(n+1+b)}{\Gamma(1+b)} \left(1 + \frac{s_1}{n+b} + \frac{s_2}{(n+b)(n+b-1)} + \dots \right). \quad (3.4)$$

It is remarkable that $b = b_1/(2b_0^2)$ and the s_i coefficients of the sub-leading $\mathcal{O}(1/n^i)$ behaviour can all be given in terms of the coefficients of the β function [7]. The relevant expressions are collected in appendix A. We also note that the scale μ_m at which m is evaluated does not appear explicitly on the right-hand side of (3.3) and hence is irrelevant in (3.1) as far as the large- n behaviour is concerned. The dependence on the scale μ of the strong coupling is compensated by the factor μ in front of $\tilde{c}_{n+1}^{(\text{as})}$ in (3.3). With these definitions the normalization N is independent of μ and μ_m . It cannot however be computed rigorously with present perturbative techniques in general, but in the limit of large negative or positive n_l it assumes the value [5]

$$\lim_{|n_l| \rightarrow \infty} N = \frac{C_F}{\pi} \times e^{\frac{5}{6}}, \quad (3.5)$$

which equals 0.97656 for $n_c = 3$ ($C_F = 4/3$).

In the following we compare the exactly known low-order coefficients of the perturbative expansion in the mass conversion relation with their expected asymptotic behaviour. By definition (see (3.3)) the normalization N is given by

$$N = \lim_{n \rightarrow \infty} \frac{c_n(\mu, \mu_m, m(\mu_m))}{c_n^{(\text{as})}(\mu, m(\mu_m))}. \quad (3.6)$$

We now determine N by evaluating the above expression for $n = 1, 2, 3, 4$, for which $c_n(\mu, \mu_m, m(\mu_m))$ is known. To this end the result of [4] for the four-loop coefficient has been expressed in terms of the strong coupling constant with n_l flavours rather than $n_l + 1$, since the asymptotic expression refers to the n_l massless flavour theory. We also use unpublished results [17] for the n_l, n_c, μ and μ_m dependence of the four-loop coefficient. In addition to the ratio $c_n/c_n^{(\text{as})}$ for n from 1 to 4 we consider the relative difference between the N estimates performed using the third and the fourth order coefficients, defined as

$$\Delta_{34} = 2 \frac{|c_3/c_3^{(\text{as})} - c_4/c_4^{(\text{as})}|}{|c_3/c_3^{(\text{as})} + c_4/c_4^{(\text{as})}|}. \quad (3.7)$$

The value of Δ_{34} can be considered to be an estimate of how close is the third order coefficient to the asymptotic value. It is likely to be an overestimate of the deviation of the

² The perturbative coefficients r_n in this reference are related to those employed here by $r_n = c_{n+1}$. With this notation the number of loops contributing to c_n is n .

fourth order coefficient from the asymptotic formula and should not be taken as an error on the normalization N .

We report our results in table 1 for $\mu_m = \overline{m}$ and the three values $\mu = \overline{m}$, $\mu = \overline{m}/2$ and $\mu = 2\overline{m}$ of the coupling renormalization scale. The number of colours has been fixed to $n_c = 3$ in this table, and the number of light flavours was varied from a very large negative value (equivalent to the large- n_l limit) up to $n_l = 10$. In columns 2 to 5 we show the ratios $c_n/c_n^{(\text{as})}$, that correspond to an estimate of N according to (3.6) for finite n . In the last column we give Δ_{34} . The \pm numbers account for the change in N due to the numerical uncertainty in the calculation of the exact four-loop conversion coefficient, which is about 0.1% on the n_l independent term for $\mu = \mu_m = m(\mu_m)$.

We first discuss the result for $\mu = \overline{m}$. For n_l very large and negative the value of N is close to the one predicted by (3.5). The value of Δ_{34} corresponds to a 4% deviation of the third order coefficient from the asymptotic result, which is indeed the case, and the fourth-order value is already much closer.³ As n_l increases, the value of N decreases, reaching 0.506(2) and 0.462(2) for $n_l = 4$ and 5, respectively, with a 9 and 13% variation when going from the third to the fourth order coefficient. As n_l increases, Δ_{34} also increases, so that for n_l above 7 the N values obtained from the third and fourth order coefficients differ by factors of order 1. This behaviour is not unexpected: by increasing the number of light flavours the first coefficient of the β function, b_0 , decreases (it vanishes for $n_l = 33/2$), hence the renormalon dominance is delayed to higher orders. We shall comment further on the n_l dependence below.

When considering different choices of the renormalization scale, we see that the $\mu = \overline{m}/2$ case leads to larger variations than $\mu = 2\overline{m}$. The large n_l limit yields a value that is about 10% higher than the exact result but the associated value $\Delta_{34} \approx 40\%$ is also large, indicating that the series is not as close to the asymptotic regime as for $\mu = \overline{m}$. For the interesting cases $n_l = 4$ and $n_l = 5$, Δ_{34} is also more than a factor of two larger than for $\mu = \overline{m}$. Again, this behaviour is not unexpected. The coefficients c_n depend only on logarithms of μ/m up to the $(n - 1)$ th power. Eq. (3.3) shows that these logarithms must asymptotically exponentiate to μ/m , which clearly happens less efficiently at finite order when $\ln(\mu/m)$ is larger. Hence we expect the best approximation to the asymptotic behaviour to occur when $\mu \approx \overline{m}$. Fig. 1 shows that this is indeed the case for large $-n_l$. It further shows a plateau around $\mu \approx \overline{m}$ and a more rapid departure from the exact result for μ smaller than \overline{m} than for larger μ , as also seen in table 1.

We also determine the normalization N for different values of n_c and show the result for Δ_{34} in fig. 2. We generically find $\Delta_{34} < 0.1$ except in regions where b_0 is small, where we do not expect our method to work. Fig. 2 therefore demonstrates that the exact four-loop coefficient indeed matches the asymptotic formula (3.1) in the expected range of n_c and n_l values, comprising those of physical interest.

For the following a reliable determination of N and an estimate of its error is particularly important for $n_c = 3$, $n_l = 5$, corresponding to the case of the top quark. We

³We may note that the contribution from sub-leading renormalon poles to c_n is of order $1/2^n$ relative to the leading one, but there is a further suppression for the case at hand due to a small numerical coefficient, at least in the large- n_l limit, see [16].

$\mu/\overline{m} = 1$					
n_l	$c_1/c_1^{(\text{as})}$	$c_2/c_2^{(\text{as})}$	$c_3/c_3^{(\text{as})}$	$c_4/c_4^{(\text{as})}$	Δ_{34}
-1000000	0.6953	0.9624	0.9349	0.9714	0.038
-10	0.4744	0.7152	0.6898	0.7005 ± 0.0002	0.015 ± 0.000
0	0.4377	0.6357	0.6130	0.5977 ± 0.0006	0.025 ± 0.001
3	0.3954	0.6150	0.5723	0.5370 ± 0.0011	0.064 ± 0.002
4	0.3633	0.6120	0.5522	0.5056 ± 0.0015	0.088 ± 0.003
5	0.3143	0.6119	0.5244	0.4616 ± 0.0020	0.127 ± 0.004
6	0.2436	0.6089	0.4818	0.3942 ± 0.0028	0.200 ± 0.007
7	0.1474	0.5378	0.4084	0.2786 ± 0.0042	0.378 ± 0.015
8	0.0098	0.0379	0.2719	0.0564 ± 0.0068	1.312 ± 0.068
10	0.2684	-0.0916	-0.1108	-1.7228 ± 0.0271	1.758 ± 0.004
$\mu/\overline{m} = 0.5$					
-1000000	1.3907	1.3554	0.6952	1.0773	0.431
-10	0.9487	0.9410	0.6701	0.7110 ± 0.0003	0.059 ± 0.000
0	0.8753	0.7907	0.6149	0.5807 ± 0.0012	0.057 ± 0.002
3	0.7908	0.7343	0.5659	0.5030 ± 0.0023	0.118 ± 0.005
4	0.7266	0.7159	0.5370	0.4631 ± 0.0030	0.148 ± 0.006
5	0.6286	0.6975	0.4943	0.4078 ± 0.0040	0.192 ± 0.010
6	0.4872	0.6704	0.4267	0.3243 ± 0.0056	0.273 ± 0.017
7	0.2948	0.5640	0.3117	0.1845 ± 0.0084	0.513 ± 0.043
8	0.0196	0.0370	0.1123	-0.0768 ± 0.0135	10.676 ± 5.676
10	0.5367	-0.0621	-0.2877	-2.1014 ± 0.0541	1.518 ± 0.011
$\mu/\overline{m} = 2$					
-1000000	0.3477	0.6235	0.8631	0.9409	0.086
-10	0.2372	0.4800	0.5883	0.6576 ± 0.0001	0.111 ± 0.000
0	0.2188	0.4380	0.5217	0.5698 ± 0.0003	0.088 ± 0.001
3	0.1977	0.4314	0.4947	0.5247 ± 0.0006	0.059 ± 0.001
4	0.1817	0.4330	0.4831	0.5026 ± 0.0007	0.040 ± 0.001
5	0.1572	0.4376	0.4681	0.4724 ± 0.0010	0.009 ± 0.002
6	0.1218	0.4413	0.4452	0.4262 ± 0.0014	0.044 ± 0.003
7	0.0737	0.3968	0.4038	0.3460 ± 0.0021	0.154 ± 0.006
8	0.0049	0.0286	0.3177	0.1877 ± 0.0034	0.515 ± 0.017
10	0.1342	-0.0761	-0.0083	-1.1238 ± 0.0135	1.971 ± 0.000

Table 1. The values of N obtained from the coefficients of the perturbative expansion up to the fourth order for several values of n_l . Three values of the renormalization scale are considered.

determine the error by varying the two renormalization scale independently, that is we vary $\mu/m(\mu_m)$ and $\mu_m/m(\mu_m)$ independently between 0.5 and 2, compute N from $c_4/c_4^{(\text{as})}$ as above, and determine the error on N from the maximal variation. The dependence of N on the two scale ratios is shown in fig. 3. With this definition our error estimate on N neither depends on the value of the heavy quark mass nor the one of the strong coupling. We find

$$N = 0.4616_{-0.070}^{+0.027}(\mu \text{ and } \mu_m) \pm 0.002(c_4). \quad (3.8)$$

As a further check we note that when the subleading term s_2 (s_1 and s_2) is removed in

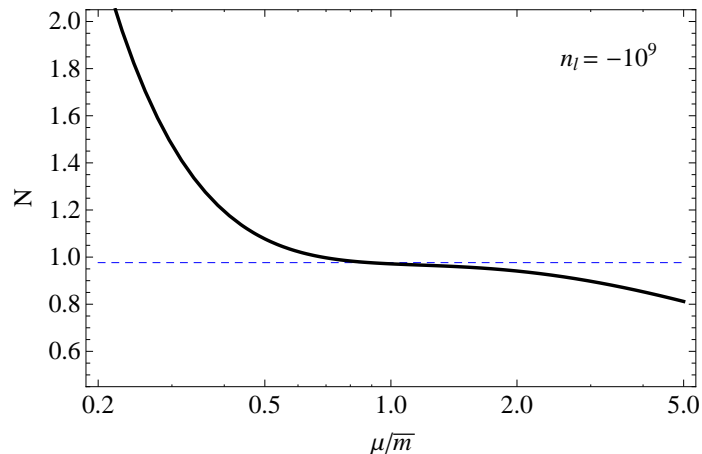


Figure 1. The normalization N as a function of μ/\bar{m} varied by a factor five around the central scale. The dashed line shows the exact value $4e^{5/6}/(3\pi) = 0.97656\dots$

(3.4), the central value changes very little to 0.4573 (0.4584).

A similar method to determine the normalization of the leading pole mass renormalon, albeit without variations of μ_m and n_c , has already been used in [12]. More precisely, instead of the four-loop pole mass considered here the three-loop static potential was employed to arrive at the best estimate, based on the fact that the pole mass and static potential leading renormalon normalizations are rigorously related by a factor of $-1/2$. Their values are indeed in good agreement with ours, though deteriorating with increasing n_l . The approach to the exact value for large negative n_l was also observed in [12].

The authors of Ref. [12] also determined the normalization N as a function of n_l and noted that it tends to zero in the range $n_l = 12 \dots 23$ close to the conformal window. We confirm this behaviour in our analysis, see Figure 4. To understand why the normalization of the leading renormalon is forced to be small in this n_l region, we look at the explicit

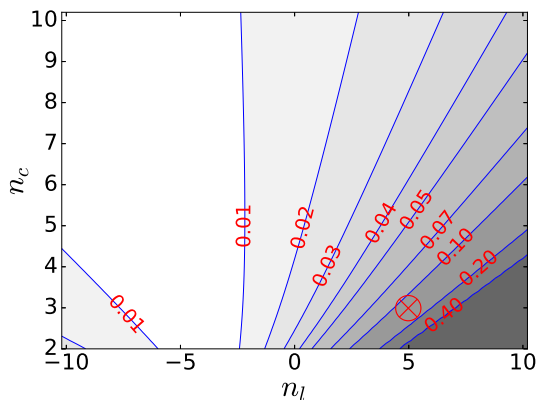


Figure 2. Δ_{34} as a function of n_c and n_l , for $\mu = \mu_m = \bar{m}$. The cross corresponds to the case relevant for top, i.e. $n_c = 3$ and $n_l = 5$.

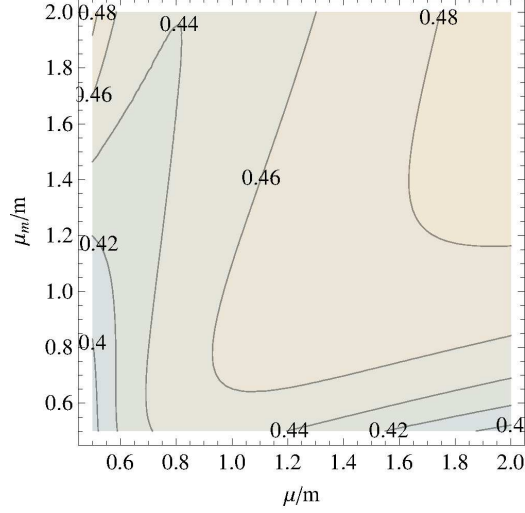


Figure 3. The normalization N as a function of $\mu/m(\mu_m)$ and $\mu_m/m(\mu_m)$.

expression of for $c_n^{(\text{as})}$ from (3.3) for $n = 4$,

$$c_4^{(\text{as})} = (2b_0)^3 (1+b)(2+b)(3+b) \left(1 + \frac{s_1}{3+b} + \frac{s_2}{(3+b)(2+b)} + \dots \right). \quad (3.9)$$

The region $n_l = 12 \dots 23$ is approximately centred around the value of n_l , where b_0 vanishes, hence $b = b_1/(2b_0^2)$ becomes large. As soon as $b \gg n_0$, where n_0 is the order from which N is determined (here $n_0 = 4$), the individual terms in the above expression behave as

$$c_4^{(\text{as})} = (2b_0)^{n_0} \left(\frac{b_1}{2b_0^2} \right)^{n_0} \left(1 + \frac{s_1}{b} + \frac{s_2}{b^2} + \dots \right) \sim \frac{1}{(2b_0)^{n_0}} \left(1 + \frac{\#}{b_0^2} + \frac{\#}{b_0^4} + \dots \right), \quad (3.10)$$

from which we conclude a) that $c_4^{(\text{as})} \sim 1/(2b_0)^{n_0}$ becomes very large, hence N must become small to fit the given value of the exact four-loop coefficient c_4 , and b) the series of sub-

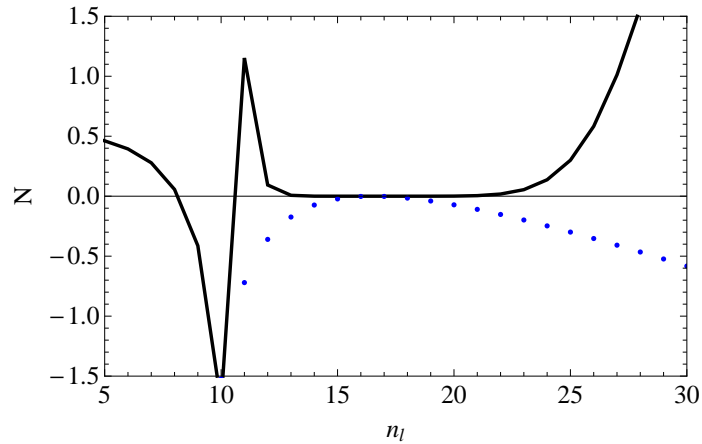


Figure 4. The normalization N (for $n_c = 3$, $\mu = \mu_m = \bar{m}$) as a function of n_l (black). The blue dots show $1/b$.

leading asymptotic terms s_1, s_2 , etc. breaks down, hence the extracted value of N is completely unreliable. The smallness of N is therefore a technical artifact of the method, which ceases to be valid when b becomes large compared to n_0 , and the question whether N is small in the conformal window cannot be answered. In fact, while small b_0 makes renormalon behaviour less relevant to low orders due to the diminished $(2b_0)^n$ factor, there seems to be no reason why the normalization N should vanish when the theory becomes conformal non-perturbatively.

4 The $m_P - \overline{m}$ conversion factor to all orders and the ultimate top pole mass uncertainty

In the following we use two methods to estimate the remainder of the mass conversion relation beyond the exactly known four-loop accuracy and to estimate the intrinsic ambiguity of summing the assumed asymptotic expansion. The first relies on truncation of the expansion and an estimate of the minimal term. The second on Borel summation. We restrict ourselves to the case of the top quark mass ($n_c = 3, n_l = 5$) and choose $\mu = \mu_m = \overline{m}$.

We begin by writing

$$m_P(n) = \overline{m} \left(1 + \sum_{k=1}^n c_k \alpha_s^k \right), \quad (4.1)$$

where the coefficients are the exact ones up to the fourth order in α_s , and determined from the asymptotic formula (3.4) (with normalization fitted to the fourth order term) for the terms of order 5 and higher. We would like to define the best value of m_P as the value at which its increment with n is minimal. More precisely, we define

$$\Delta(n + 1/2) = m_P(n + 1) - m_P(n), \quad (4.2)$$

which is a decreasing function of n up to a certain value n_0 beyond which it begins to increase due to the renormalon divergence of the series expansion. By interpolating Δ with a quadratic form in the three points $n_0 - 1/2, n_0 + 1/2, n_0 + 3/2$, we find its minimum at (generally non-integer)

$$n_{\min} = n_0 + 1/2 - \frac{\Delta(n_0 + 3/2) - \Delta(n_0 - 1/2)}{2(\Delta(n_0 + 3/2) + \Delta(n_0 - 1/2) - 2\Delta(n_0 + 1/2))}. \quad (4.3)$$

By interpolating linearly the value of $m_P(n_{\min})$ between n_0 and $n_0 + 1$ we get

$$m_P^c = \frac{m_P(n_0)(\Delta(n_0 + 3/2) - \Delta(n_0 + 1/2)) + m_P(n_0 + 1)(\Delta(n_0 - 1/2) - \Delta(n_0 + 1/2))}{\Delta(n_0 + 3/2) + \Delta(n_0 - 1/2) - 2\Delta(n_0 + 1/2)} \quad (4.4)$$

as the best value of the pole mass. We note that with this prescription, if $\Delta(n_0 - 1/2) = \Delta(n_0 + 3/2)$, then m_P^c corresponds to $(m_P(n_0) + m_P(n_0 + 1))/2$, as one would intuitively expect, while for $\Delta(n_0 - 1/2) \gg \Delta(n_0 + 3/2)$ ($\Delta(n_0 - 1/2) \ll \Delta(n_0 + 3/2)$), we obtain $m_P(n_0 + 1)$ ($m_P(n_0)$).

We now estimate the correction to the top pole mass due to terms of order higher than four by

$$\delta^{(5+)} m_P = N \mu \sum_{k=5}^{\infty} \tilde{c}_k^{(\text{as})} \alpha_s^k(\mu), \quad (4.5)$$

j	$\tilde{c}_j^{(\text{as})}$	$\tilde{c}_j^{(\text{as})} \alpha_s^j$
5	0.985499×10^2	0.001484
6	0.641788×10^3	0.001049
7	0.495994×10^4	0.000880
8	0.443735×10^5	0.000854
9	0.451072×10^6	0.000942
10	0.513535×10^7	0.001164
11	0.647283×10^8	0.001593
12	0.894824×10^9	0.002390
13	0.134620×10^{11}	0.003902
14	0.218949×10^{12}	0.006888
15	0.382818×10^{13}	0.013070

Table 2. The coefficients $\tilde{c}_j^{(\text{as})}$ above the fourth order. Their value multiplied by the corresponding power of $\alpha_s = 0.108531$ is also reported.

where $\tilde{c}_j^{(\text{as})}$ is defined in (3.4), and the barred sum represents the procedure we have just outlined for the evaluation of the (divergent) sum. We report in table 2 the values of $\tilde{c}_j^{(\text{as})}$ beyond the fourth order term. Eq. (4.5) can be easily computed for any value of α_s and μ and is well approximated by the second-order Taylor series around the reference value:

$$\delta^{(5+)} m_P = N \mu \times 10^{-3} \left(3.604 + 14.69 \left(\frac{\alpha_s(\mu)}{0.1085} - 1 \right) + 9.54 \left(\frac{\alpha_s(\mu)}{0.1085} - 1 \right)^2 \right). \quad (4.6)$$

For typical values of $N \approx 0.5$ and $\mu \approx 160$ GeV the formula is accurate at the sub-MeV level for a $\pm 5\%$ variation of the strong coupling constant.

We now adopt the PDG value $\alpha_s(M_Z) = 0.1181 \pm 0.0013$, and take $\mu = \bar{m} = 163.508$ GeV for definiteness. With this input we find $\alpha_s(\mu) = 0.108531$ for the (five flavour) strong coupling constant and 173.34 GeV for the top pole mass using the four-loop conversion formula. From the values reported in the table and the value of N given in (3.8) we obtain for the series remainder

$$\delta^{(5+)} m_P = 0.272_{-0.041}^{+0.016} (N) \pm 0.001 (c_4) \pm 0.011 (\alpha_s) \pm 0.066 (\text{ambiguity}) \text{ GeV}, \quad (4.7)$$

where we show the error due to the uncertainty in the normalization N , the four-loop coefficient c_4 , and $\alpha_s(M_Z)$. For the irreducible renormalon ambiguity we tentatively estimate the size of the first omitted term by the value of $\Delta(n_0 - 1/2)$. For the top mass conversion factor we find

$$m_P^c / \bar{m} = 1.06177_{-0.00025}^{+0.00010} (N) \pm 0.00001 (c_4) \pm 0.00087 (\alpha_s) \pm 0.00041 (\text{ambiguity}). \quad (4.8)$$

We also computed the change of the conversion factor under variations of μ/\bar{m} and μ_m/\bar{m} , simultaneously in the exact four-loop part and the remainder, accounting for the dependence

of N on μ and μ_m (fig. 3). This leads to ${}_{-0.00041}^{+0.00025}$, which we do not include above, since it is strongly correlated with the uncertainty of the same order from N alone.

In the second method we first compute the Borel transform of the asymptotic series coefficients $\tilde{c}^{(\text{as})}$ in (3.3), which gives

$$B[\tilde{c}^{(\text{as})}](t) = \frac{1}{(1-2b_0t)^{1+b}} + \frac{s_1}{b} \frac{1}{(1-2b_0t)^b} + \frac{s_2}{b(b-1)} \frac{1}{(1-2b_0t)^{-1+b}} + \dots, \quad (4.9)$$

and then the Borel sum

$$BS[\tilde{c}^{(\text{as})}](\alpha_s) = \int_0^\infty dt e^{-t/\alpha_s} B[\tilde{c}^{(\text{as})}](t). \quad (4.10)$$

Since the series is not Borel-summable due to the IR renormalon singularity at $t = 1/(2b_0)$, we define the sum as the principal value and estimate the ambiguity as the imaginary part of the integral when the contour is deformed into the upper complex plane, divided by Pi. This procedure is known to usually give a reliable estimate [16], close to the sum to the minimal term and the estimate of the summation ambiguity by the smallest term in the series. The Borel sum can easily be computed analytically, since (with the contour deformed into the upper complex plane)

$$\int_0^\infty dt e^{-t/\alpha_s} \frac{1}{(1-2b_0t)^\gamma} = \frac{\alpha_s}{(-2b_0\alpha_s)^\gamma} e^{-1/(2b_0\alpha_s)} \Gamma(1-\gamma, -1/(2b_0\alpha_s)), \quad (4.11)$$

where $\Gamma(a, z)$ denotes the incomplete Gamma function. The remainder of the mass conversion formula is obtained by subtracting the first four coefficients, resulting in

$$\delta^{(5+)} m_P = N\mu \left(BS[\tilde{c}^{(\text{as})}](\alpha_s(\mu)) - \sum_{k=1}^4 \tilde{c}_k^{(\text{as})} \alpha_s(\mu)^k \right). \quad (4.12)$$

With parameter input as above, we find

$$\delta^{(5+)} m_P = 0.250_{-0.038}^{+0.015} (N) \pm 0.001 (c_4) \pm 0.010 (\alpha_s) \pm 0.071 (\text{ambiguity}) \text{ GeV}, \quad (4.13)$$

which is close to the result (4.7) from the previous method. For any value of α_s and μ the result can again be determined accurately in the phenomenologically relevant region according to the fit formula

$$\delta^{(5+)} m_P = N\mu \times 10^{-3} \left(3.315 + 12.71 \left(\frac{\alpha_s(\mu)}{0.1085} - 1 \right) + 4.55 \left(\frac{\alpha_s(\mu)}{0.1085} - 1 \right)^2 \right). \quad (4.14)$$

For the top mass conversion factor itself, we find

$$m_P^c / \overline{m} = 1.06164_{-0.00023}^{+0.00009} (N) \pm 0.00001 (c_4) \pm 0.00086 (\alpha_s) \pm 0.00043 (\text{ambiguity}). \quad (4.15)$$

In this case, the scale variation is ${}_{-0.00028}^{+0.00013}$. We adopt (4.13) and (4.15) as our final results. Given the $\overline{\text{MS}}$ mass, the top quark pole mass is determined by this relation with an accuracy of 0.92 per mil, half of which is due to the irreducible uncertainty of the relation itself.

5 Conclusions

We employed the four-loop coefficient in the pole- $\overline{\text{MS}}$ quark mass relation, which has recently become available [4], and knowledge of the leading asymptotic behaviour of the series expansion of the mass conversion factor [7] to estimate the remainder of the series from terms above the four-loop order and the intrinsic ambiguity due to the asymptotic nature of the series. For the case of the top quark we find about 250 MeV for the former and 70 MeV for the ambiguity, which also represents the ultimate precision that can be obtained for the pole mass. The ambiguity of 70 MeV is far below the accuracy that can conceivably be achieved at the Large Hadron Collider, but larger than the one foreseen in theoretical and experimental studies [18, 19] of a scan of the top pair production threshold at a high-energy e^+e^- collider. In this case the pole mass ceases to be a useful concept and other mass definitions must be employed.

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A Summary of formulae

In this Appendix, in order to make contact with the notation of [7, 16], we define the QCD beta-function as

$$\beta(\alpha_s) = \mu^2 \frac{\partial \alpha_s(\mu)}{\partial \mu^2} = \beta_0 \alpha_s^2 + \beta_1 \alpha_s^3 + \dots, \quad (\text{A.1})$$

With this convention $\beta_0 = -(11n_c/3 - 2n_l/3)/(4\pi)$, while in the main text we used $b_i = -\beta_i > 0$ (for small n_l). We adopt the $\overline{\text{MS}}$ scheme with n_l massless quark flavours. (The heavy quark whose mass is considered here is decoupled.) The constants that appear in (3.4) are given by [7, 16] $b = -\beta_1/(2\beta_0^2)$ and

$$s_1 = \left(-\frac{1}{2\beta_0}\right) \left(-\frac{\beta_1^2}{2\beta_0^3} + \frac{\beta_2}{2\beta_0^2}\right), \quad (\text{A.2})$$

$$s_2 = \left(-\frac{1}{2\beta_0}\right)^2 \left(\frac{\beta_1^4}{8\beta_0^6} + \frac{\beta_1^3}{4\beta_0^4} - \frac{\beta_1^2\beta_2}{4\beta_0^5} - \frac{\beta_1\beta_2}{2\beta_0^3} + \frac{\beta_2^2}{8\beta_0^4} + \frac{\beta_3}{4\beta_0^2}\right). \quad (\text{A.3})$$

Note that we have corrected some misprints in the expression for b and s_2 given in [16] (eqs. (5.91) and (5.92)) as already noted in [10].

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