

Single-differential top quark pair production cross sections with running mass schemes at NLO

Toni Mäkelä,^{a,*} André Hoang,^b Katerina Lipka^{a,c} and Sven-Olaf Moch^d

^a*Deutsches Elektronen-Synchrotron,
Notkestraße 85, Hamburg, Germany*

^b*Fakultät für Physik, Universität Wien,
Boltzmannngasse 5, Wien, Austria*

^c*Fakultät für Mathematik und Naturwissenschaften, Bergische Universität Wuppertal,
Gaussstrasse 20, Wuppertal, Germany*

^d*II. Institut für Theoretische Physik, Universität Hamburg,
Luruper Chaussee 149, Hamburg, Germany*

E-mail: toni.maekela@desy.de

Single-differential cross section predictions for top quark pair production are presented at next-to-leading order, using running top quark mass renormalization schemes. The evolution of the mass of the top quark is performed in the MSR scheme $m_t^{\text{MSR}}(\mu)$ for renormalization scales μ below the $\overline{\text{MS}}$ top quark mass $\overline{m}_t(\overline{m}_t)$, and in the $\overline{\text{MS}}$ scheme $\overline{m}_t(\mu)$ for scales above. In particular, the implementation of a mass renormalization scale independent of the strong coupling renormalization scale and factorization scale in quantum chromodynamics allows investigating independent dynamical scale variations. Furthermore, the first theoretically consistent extraction of the top quark MSR mass from experimental data is presented.

*41st International Conference on High Energy physics - ICHEP2022
6-13 July, 2022
Bologna, Italy*

*Speaker

The top quark mass m_t is a fundamental parameter of the Standard Model and has an important role in many predictions, both directly and via higher-order corrections. Yet, the formal definition of quark masses makes them renormalization scheme dependent quantities. While the pole mass m_t^{pole} suffers from the renormalon ambiguity, an infrared sensitivity of the order of the scale of quantum chromodynamics (QCD) [1], short-distance masses e.g. the $\overline{\text{MS}}$ mass $\overline{m}_t(\mu_m)$ and the MSR mass $m_t^{\text{MSR}}(R)$ [2] do not. However, the dependence on the mass renormalization scales μ_m and R necessitates proper scale setting for the extraction of theoretically well-defined masses from cross section measurements to avoid the appearance of large logarithms.

The pole and $\overline{\text{MS}}$ masses are related by $m_t^{\text{pole}} = \overline{m}_t(\mu_m) \left(1 + \sum_{n=1}^{\infty} d_n^{\overline{\text{MS}}}(\mu_m) (a_S(\mu_m))^n \right)$, where $a_S \equiv \alpha_S/\pi$, and $d_n^{\overline{\text{MS}}}(\mu_m)$ are perturbative coefficients. The pole and MSR mass relation reads $m_t^{\text{pole}} = m_t^{\text{MSR}}(R) + R \sum_{n=1}^{\infty} d_n^{\text{MSR}}(a_S(R))^n$, i.e. $m_t^{\text{MSR}}(R)$ approaches m_t^{pole} in the formal limit $R \rightarrow 0$, and the $\overline{\text{MS}}$ mass at $R \rightarrow \overline{m}_t(\overline{m}_t)$ up to a small matching correction. The latter are obtained by integrating out top quark loop corrections at $R \lesssim \overline{m}_t(\overline{m}_t)$ [3]. The R -evolution of $m_t^{\text{MSR}}(R)$ is linear, contrary to the logarithmic μ_m evolution of $\overline{m}_t(\mu_m)$, and designed to capture the correct physical logarithms for observables with m_t dependence generated at dynamical scales $R < m_t$ (e.g. resonances, thresholds, low-energy endpoints) [4]. For dynamical scales of order and larger than m_t , the $\overline{\text{MS}}$ mass and evolution are used. Based on Ref. [5], the top quark-antiquark ($t\bar{t}$) production cross section as a function of the $t\bar{t}$ system invariant mass $m_{t\bar{t}}$ at next-to-leading order (NLO) reads

$$\frac{d\sigma}{dm_{t\bar{t}}} = (a_S)^2 \frac{d\sigma^{(0)}}{dm_{t\bar{t}}}(m, \mu_r, \mu_f) + (a_S)^3 \frac{d\sigma^{(1)}}{dm_{t\bar{t}}}(m, \mu_r, \mu_f) + (a_S)^3 \tilde{R} d_1 \frac{d}{dm_t} \left(\frac{d\sigma^{(0)}(m, \mu_r, \mu_f)}{dm_{t\bar{t}}} \right) \Big|_{m_t=m}, \quad (1)$$

where $\sigma^{(0)}$ is the leading order (LO) and $\sigma^{(1)}$ the NLO cross section in the pole mass scheme and the derivative term at NLO implements the $\overline{\text{MS}}$ or MSR top mass schemes, the renormalization (factorization) scale μ_r (μ_f) is independent of the mass renormalization scales R or μ_m and $a_S = a_S(\mu_r)$. Furthermore,

$$(m, d_1, \tilde{R}) = \begin{cases} (m_t^{\text{MSR}}(R), d_1^{\text{MSR}}, R), & \text{in the MSR regime } (R < \overline{m}_t(\overline{m}_t)), \\ (\overline{m}_t(\mu_m), d_1^{\overline{\text{MS}}}(\mu_m), \overline{m}_t(\mu_m)), & \text{in the } \overline{\text{MS}} \text{ regime } (R > \overline{m}_t(\overline{m}_t)). \end{cases} \quad (2)$$

In this work, the cross section given in Eq. (1) is implemented into MCFM v6.8 [6]. The running of m_t and the $t\bar{t}$ cross section as a function of $m_{t\bar{t}}$ are shown in Fig. 1.

The implementation of the mass renormalization scales independently from μ_r and μ_f allows the first investigation of the dependence on the scale R . As illustrated in Fig. 2 for $d\sigma/dm_{t\bar{t}}$ cross section in the bin $m_{t\bar{t}} \in [333, 366]$ GeV, low values of μ_r and μ_f result in quick stabilization of the NLO $t\bar{t}$ cross section as a function of the mass renormalization scale in this bin, which contains high sensitivity to m_t . Furthermore, Fig. 2 indicates unsmooth behavior at low R , and the cross section stabilizes at $R \gtrsim 60$ GeV. This is attributed to Coulomb effects spoiling the convergence of expansions in a_S in fixed-order QCD near the $t\bar{t}$ production threshold. In future studies, this is to be solved by including quasi-boundstate corrections and the resummation of soft gluon effects, following e.g. Ref. [7]. In accordance with these observations, the central values of R , μ_r and μ_f are set to 80 GeV in the following to obtain predictions robust against scale variations.

A determination of $m_t^{\text{MSR}}(R)$ is performed using the single-differential $t\bar{t}$ production cross section measured by the CMS Collaboration in pp collisions at $\sqrt{s} = 13$ TeV [8], corresponding to

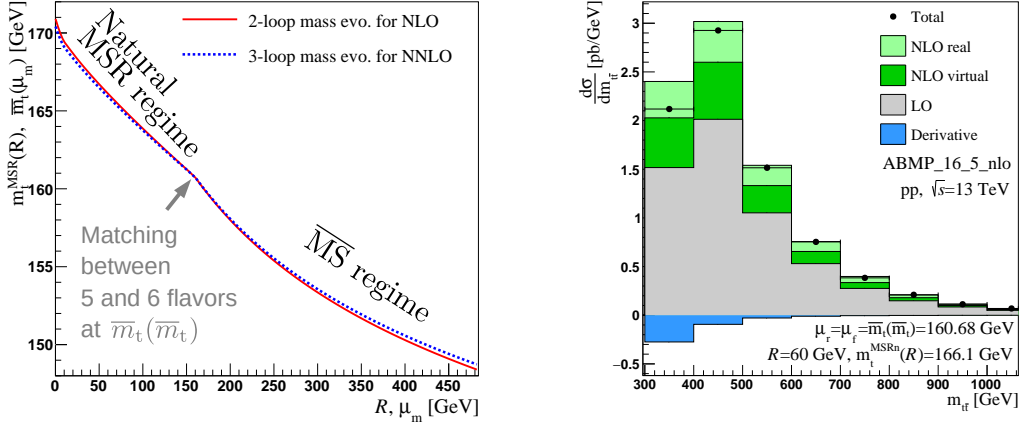


Figure 1: *Left:* the running of m_t in the MSR and $\overline{\text{MS}}$ evolution regimes, see also Ref. [2]. *Right:* the $t\bar{t}$ production cross section at NLO in bins of $m_{t\bar{t}}$ (dots) and the contributions of the terms in Eq. (1) (histograms).

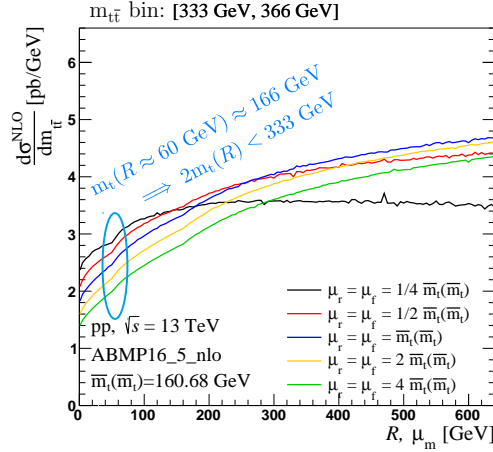


Figure 2: The $m_{t\bar{t}} \in [333, 366]$ GeV bin of $d\sigma/dm_{t\bar{t}}$ as a function of R, μ_m . At $R \lesssim 60$ GeV, threshold effects are prominent. Low values of μ_r, μ_f are observed to stabilize the predictions as a function of R, μ_m .

an integrated luminosity of 35.9 fb^{-1} . The cross section is provided in four bins: $m_{t\bar{t}} < 420$ GeV, $m_{t\bar{t}} \in [420, 550]$ GeV, $m_{t\bar{t}} \in [550, 810]$ GeV and $m_{t\bar{t}} > 810$ GeV. The top quark MSR mass is extracted by fitting $t\bar{t}$ production cross section predictions, computed with the ABMP16 5 flavor PDF [9] at NLO, to the experimental data. For the extraction, $R = 80$ GeV is assumed. For comparison with previous studies, the resulting $m_t^{\text{MSR}}(80 \text{ GeV})$ is evolved to the reference scale $R = 1$ GeV, as well as translated to $\bar{m}_t(\bar{m}_t)$. The fit uncertainty is obtained via the $\Delta\chi^2 = 1$ tolerance criterion. The uncertainty in the initial R choice is estimated by repeating the fits at $R = 60$ GeV and 100 GeV, and taking the difference of the masses evolved to the reference scales to the respective results of the $R = 80$ GeV fit. The μ_r, μ_f uncertainty is obtained by multiplying the scales independently by $2^{\pm 1}$, avoiding cases where one scale is multiplied by 2 and the other by $1/2$, and constructing an envelope.

With $\mu_r = \mu_f = m_t^{\text{MSR}}(80 \text{ GeV})$ throughout the $m_{t\bar{t}}$ distribution, evolving the obtained

$m_t^{\text{MSR}}(80 \text{ GeV})$ to $R = 1 \text{ GeV}$ yields $m_t^{\text{MSR}}(1 \text{ GeV}) = 173.2 \pm 0.6 (\text{fit})_{-0.6}^{+0.4} (\mu_r, \mu_f)_{-0.5}^{+0.4} (R) \text{ GeV}$, compatible with the Monte Carlo calibration studies performed in Ref. [10] where $R = 1 \text{ GeV}$ was also adopted as the reference scale. Furthermore, the result translates into $\bar{m}_t(\bar{m}_t) = 163.3_{-1.0}^{+0.8} \text{ GeV}$, which is in agreement with Ref. [8]. However, it disagrees with Ref. [11], where $m_t^{\text{pole}} = 170.5 \pm 0.8 \text{ GeV}$ was obtained, which translates into $m_t^{\text{MSR}}(1 \text{ GeV}) = 170.2 \pm 0.8 \text{ GeV}$ interpreting the pole mass [11] as the asymptotic pole mass [4].

Setting μ_r and μ_f to $m_t^{\text{MSR}}(80 \text{ GeV})/2$ for $m_{\bar{t}} < 420 \text{ GeV}$ and to $m_t^{\text{MSR}}(80 \text{ GeV})$ for $m_{\bar{t}} > 420 \text{ GeV}$ yields $m_t^{\text{MSR}}(1 \text{ GeV}) = 174.8 \pm 0.5 (\text{fit})_{-0.4}^{+0.2} (\mu_r, \mu_f)_{-0.3}^{+0.2} (R) \text{ GeV}$. As expected from the present investigations, the setting increases robustness against scale variations, resulting in small uncertainties. Though a full understanding of m_t extracted from cross section measurements requires the inclusion of threshold Coulomb effects, the results indicate that the choice of the top quark mass scheme and the use of dynamical renormalization scales have a considerable impact on the phenomenological analysis and need to be investigated thoroughly.

References

- [1] I. I. Y. Bigi, M. A. Shifman, N. G. Uraltsev and A. I. Vainshtein, Phys. Rev. D **50** (1994), 2234-2246, arXiv:hep-ph/9402360 [hep-ph]; M. Beneke and V. M. Braun, Nucl. Phys. B **426** (1994), 301-343, arXiv:hep-ph/9402364 [hep-ph]; M. C. Smith and S. S. Willenbrock, Phys. Rev. Lett. **79** (1997), 3825-3828, arXiv:hep-ph/9612329 [hep-ph].
- [2] A. H. Hoang, C. Lepenik and M. Preisser, JHEP **09** (2017), 099, arXiv:1706.08526 [hep-ph]; A. H. Hoang et al., Phys. Rev. Lett. **101** (2008), 151602, arXiv:0803.4214 [hep-ph].
- [3] A. H. Hoang, A. Jain, C. Lepenik, V. Mateu, M. Preisser, I. Scimemi and I. W. Stewart, JHEP **04** (2018), 003, arXiv:1704.01580 [hep-ph].
- [4] A. H. Hoang, C. Lepenik and V. Mateu, Comput. Phys. Commun. **270** (2022), 108145, arXiv:2102.01085 [hep-ph].
- [5] M. Dowling and S. O. Moch, Eur. Phys. J. C **74** (2014) no.11, 3167, arXiv:1305.6422 [hep-ph].
- [6] J. M. Campbell and R. K. Ellis, Nucl. Phys. B Proc. Suppl. **205-206** (2010), 10-15, arXiv:1007.3492 [hep-ph]; J. M. Campbell and R. K. Ellis, J. Phys. G **42** (2015) no.1, 015005, arXiv:1204.1513 [hep-ph].
- [7] Y. Kiyo, J. H. Kuhn, S. Moch, M. Steinhauser and P. Uwer, Eur. Phys. J. C **60** (2009), 375-386, arXiv:0812.0919 [hep-ph].
- [8] CMS Collaboration, Phys. Lett. B **803** (2020), 135263, arXiv:1909.09193 [hep-ex].
- [9] S. Alekhin, J. Blümlein, S. Moch and R. Placakyte, Phys. Rev. D **96** (2017) no.1, 014011, arXiv:1701.05838 [hep-ph].
- [10] ATLAS Collaboration, ATL-PHYS-PUB-2021-034.
- [11] CMS Collaboration, Eur. Phys. J. C **80** (2020) no.7, 658, arXiv:1904.05237 [hep-ex].