Selected Topics in Top Quark Physics

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Abstract

I discuss three different selected topics in top quark physics. The first topic concerns the next-to-next-to-leading order calculation of the hadroproduction of top quark pairs and the role of multiple polylogarithms in this calculation. I report on an ongoing next-to-next-to-leading order calculation of heavy quark pair production in hadron collisions where the loop–by–loop part of the calculation is about to be completed. Calculating the loop-by-loop part allows one to take a glimpse at the mathematical structure of the full NNLO calculation. The loop-by-loop contributions bring in a new class of functions introduced only eight years ago by the Russian mathematician Goncharov called multiple polylogarithms. The second topic concerns a next-to-leading order calculation of unpolarized top quark decays which are analyzed in cascade fashion $t \rightarrow b + W^+$ followed by $W^+ \rightarrow l^+ + \nu_l$. Finally, I present some next-to-leading order results on polarized top quark decays which are analyzed in the top quark rest system.
1 Introduction

I begin my talk with a few remarks on present top quark yields at the Tevatron and on expected top quark yields at the LHC which will start running at the end of 2007. After a slow start in early 2001 Tevatron II started reaching design peak luminosities of \(8.5 \times 10^{31} \text{cm}^{-2} \text{s}^{-1}\) in 2004. The best weekly performance was early 2006 with a weekly integrated luminosity of 25 pb\(^{-1}\). If Tevatron II could perform at this rate it would be able to collect 1.3 fb\(^{-1}\) in a year. There has been a three months shutdown in the spring of 2006 with some (electron cooling) improvements on the \(\bar{p}\) beam. The hope was that there will be a factor two or three improvement in luminosity after the shutdown. Such a factor would be dearly needed if one wants to reach the projected total of 8 fb\(^{-1}\) when the machine is closed down in 2009. At the time of writing this factor has not been realized so far after a few months of post–shutdown running although the machine is performing quite well with continuous improvements. At a c.m. energy of \(\sqrt{s} = 1.96\text{TeV}\) with \(\sigma(t\bar{t}) \approx 6.8\text{pb}\), one expects around 7000 \(t\bar{t}\) pairs at each detector (CDF and DO) for an integrated luminosity of 1 fb\(^{-1}\). Single top production occurs at about 33% of the \(t\bar{t}\) pair production rate but has not been detected so far.

Much bigger samples of top quarks will be available at the LHC. Due to the higher energy of the LHC the cross section increases by a factor of 100. Also there will be a ten–fold increase in luminosity at the LHC. Thus one will have \(10^7\) \(t\bar{t}\)–pairs per year, or one \(t\bar{t}\)–pair every four seconds at each detector (ATLAS and CMS). In a later high luminosity run there will be another factor of ten increase in luminosity such that one will have a \(t\bar{t}\)–pair produced every half second. Again single top production occurs at approximately one–third the rate of \(t\bar{t}\)–production. Singly produced top quarks will be highly polarized because they are produced weakly. This opens the way to study angular correlations between the polarization of the top quark and its decay products which forms the third topic of this talk.

The yield of top quark pairs at the International Linear Collider (possibly starting in 2015) will be \(\approx (1–4)\cdot 10^5/\text{y}\) depending on the c.m. energy \(\approx 360–800\text{GeV}\). In \(e^+–e^-\) interactions a high degree of polarization of the top (or antitop) quark can be achieved through tuning of the beam polarization.

2 NNLO description of heavy top quark production

The full next-to-leading order (NLO) QCD corrections to hadroproduction of heavy flavors were completed as early as 1988 [1, 2]. They have raised the leading order (LO) estimates [3] but were still below the experimental results on bottom quark pair production (see e.g. [4]). In a recent analysis theory moved closer to experiment [5]. First experimental results on hadronic \(t\bar{t}\)–pair production [5, 6] are in agreement with theoretical NLO QCD predictions [7, 8] within the large theoretical and experimental error bars.

A large uncertainty in the NLO calculation results from the freedom in the choice of the renormalization and factorization scales. These scale uncertainties amount to a \(\approx 10\%\) theoretical error in the NLO cross section predictions [2]. The dependence on the factorization and renormalization scales is expected to be greatly reduced at next-to-next-to-leading order (NNLO). This will reduce the theoretical uncertainty.
In Fig. 1 we show one generic diagram each for the four classes of gluon induced contributions that need to be calculated for the NNLO corrections to hadroproduction of heavy flavors. They involve the two-loop contribution (Fig. 1a), the loop-by-loop contribution (Fig. 1b), the one-loop gluon emission contribution (Fig. 1c) and, finally, the two gluon emission contribution (Fig. 1d). In our work we have concentrated on the loop-by-loop contributions exemplified by Fig. 1b. Specifically, working in the framework of dimensional regularization, we have calculated $O(\epsilon^2)$ results for all scalar massive one-loop one-, two-, three- and four-point integrals that are needed in the calculation of hadronic heavy flavour production [10]. Because the one-loop integrals exhibit infrared (IR)/collinear (M) singularities up to $O(\epsilon^{-2})$ [9] one needs to know the one-loop integrals up to $O(\epsilon^2)$ since the one-loop contributions appear in product form in the loop-by-loop contributions. It is exactly the $O(\epsilon^2)$ terms in the scalar massive three- and four-point integrals that bring in multiple polylogarithms [10, 11].

Calculating the loop-by-loop contributions allows one to obtain a glimpse of the complexity that is waiting for us in the full NNLO calculation. This complexity does in fact reveal itself in terms of a very rich polylogarithmic structure of the Laurent series expansion of the scalar one-loop integrals as well as the appearance of multiple polylogarithms of maximal weight and depth four.

To underscore the statement that the loop–by–loop contributions reveal the mathematical structure of a full NNLO calculation let us take a look at the paper by Bernreuther et al. [12] who calculated the $O(\epsilon^1)$ contributions to the one–loop vertex correction of the process $V \rightarrow Q\bar{Q}$. These are needed for the loop–by–loop part of a NNLO calculation of heavy quark pair production in $e^+e^-$–annihilations. The result can be expressed in terms of one–dimensional harmonic polylogarithms of maximal weight three. The same paper
also lists results on the corresponding $O(\epsilon^0)$ two-loop vertex corrections which contain one-dimensional harmonic polylogarithms of maximal weight four. This shows that the same mathematical complexity appears in the loop–by–loop contributions as in the two-loop contribution. Let me mention that heavy quark pair production in $e^+ - e^-$–annihilations is a somewhat simpler problem than heavy quark pair production in hadronic collisions because of the appearance of one additional mass scale in the latter case. This explains why one has only one-dimensional harmonic polylogarithms in $e^+ - e^-$–annihilations case compared to the multiple polylogarithms appearing in the hadronic collision calculation. Even then the two–loop vertex correction to $V \rightarrow Q\bar{Q}$ listed in [12] takes up more than twelve pages.

The scalar four–point integrals appearing in the calculation of the loop–by–loop evaluation are the most difficult to calculate. They contain a very rich structure in terms of polylogarithmic functions. For example, the $\epsilon^2$–coefficients of the Laurent series expansion of the four–point integrals contain logarithms and classical polylogarithms up to order four (i.e. $\text{Li}_4$) in conjunction with the $\zeta$–functions $\zeta(2), \zeta(3)$ and $\zeta(4)$ and products thereof, and a new class of functions which are now termed multiple polylogarithms [13].

Since this is a conference on mathematical physics it is appropriate to dwell a little on the subject of multiple polylogarithms. A multiple polylogarithm is represented by

\[
\text{Li}_{m_k,\ldots,m_1}(x_k,\ldots,x_1) = \int_0^{x_1 x_2 \ldots x_k} \left( \frac{dt}{t} \right)^{m_1-1} \int_0^{x_2 x_3 \ldots x_k} \left( \frac{dt}{t} \right)^{m_2-1} \ldots \int_0^{x_k} \left( \frac{dt}{t} \right)^{m_k-1} \frac{dt}{1-t},
\]

where the iterated integrals are defined by

\[
\int_0^{t_n} \frac{dt}{a_n-t} \circ \ldots \circ \frac{dt}{a_1-t} = \int_0^{t_n} \frac{dt}{a_n-t_n} \int_0^{t_{n-1}} \frac{dt}{a_{n-1}-t_{n-1}} \times \ldots \times \int_0^{t_2} \frac{dt}{a_1-t_1}.
\]

The indices $m_k$ and $k$ are positive integers. The multiple polylogarithms are classified according to their weight $w = m_1 + m_2 + \ldots + m_k$ and their depth $k$. We mention that a very efficient program for the numerical evaluation of multiple polylogarithms has recently been developed in Mainz which, characteristically, is based on the language GiNaC [14].

The classical polylogarithms, Nielsen’s generalized polylogarithms, the one– and two–dimensional harmonic polylogarithms are all special cases of Goncharov’s multiple polylogarithms (see e.g. [15]). For example, the classical polylogarithms

\[
\text{Li}_n(z) = \int_0^z \frac{\text{Li}_{n-1}(x)}{x} dx \quad n \geq 2; \quad \text{Li}_1(z) = -\ln(1-z) \quad (1)
\]

are multiple polylogarithms of weight $n$ and depth 1.
In our original Feynman parameter calculation our results were written down as one-dimensional integral representations given by the integrals

\[ F_{\sigma_1\sigma_2\sigma_3}(\alpha_1, \alpha_2, \alpha_3, \alpha_4) = \int_0^1 dy \frac{\ln(\alpha_1 + \sigma_1 y) \ln(\alpha_2 + \sigma_2 y) \ln(\alpha_3 + \sigma_3 y)}{\alpha_4 + y} \] (2)

and

\[ F_{\sigma_1}(\alpha_1, \alpha_2, \alpha_3, \alpha_4) = \int_0^1 dy \frac{\ln(\alpha_1 + \sigma_1 y) \operatorname{Li}_2(\alpha_2 + \alpha_3 y)}{\alpha_4 + y} , \] (3)

where the \( \sigma_i \) take values \( \pm 1 \) and the \( \alpha_j \)'s are combinations of the kinematical variables of the process. The numerical evaluation of these one-dimensional integral representations are quite stable. The functions \( F_{\sigma_1\sigma_2\sigma_3}(\alpha_1, \alpha_2, \alpha_3, \alpha_4) \) and \( F_{\sigma_1}(\alpha_1, \alpha_2, \alpha_3, \alpha_4) \) are related to multiple polylogarithms of maximal weight and depth four as shown in [11].

We are now in the process of computing the full loop–by–loop contributions including the spin and colour algebra arising from squaring the full one–loop amplitudes as given in [16]. A first result has been obtained for the Abelian case of photon–photon production of heavy quark pairs [17].

### 3 Decays of unpolarized and polarized top quarks

After this brief mathematical detour I return to the physics of top quark decays. In the SM the top quark decays almost 100\% to a \( W^+ \) and a bottom quark. Also, the top quark decays so fast that it retains its initial polarization when it decays. I describe both unpolarized and polarized top quark decays. In the unpolarized case I analyze top quark decays in cascade fashion as a two step process involving the decay \( t \rightarrow b + W^+ \) and \( W^+ \rightarrow l^+ + \nu_l \) in the respective rest frames of the top quark and the \( W^+–\)boson. In the polarized case I perform the decay analysis in the rest system of the top quark. I discuss polar and azimuthal correlations involving the polarization of the top quark and the momenta of the decay products in the decay \( t(\uparrow) \rightarrow X_b + l^+ + \nu_l \).

The decay of a polarized top quark into a \( W^+–\)boson and a jet with \( b \)–quantum numbers \( t(\uparrow) \rightarrow X_b + l^+ + \nu_l \) is described by altogether eight invariant structure functions (see e.g. [18] [19]).

\[
H^{\mu\nu} = \left(-g^{\mu\nu} H_1 + p_t^\mu p_t^\nu H_2 - i\epsilon^{\mu\nu\rho\sigma} p_{t\rho} q_{\sigma} H_3 \right) + \right.

\left[ (q \cdot s_t) \left(-g^{\mu\nu} G_1 + p_t^\mu p_t^\nu G_2 - i\epsilon^{\mu\nu\rho\sigma} p_{t\rho} q_{\sigma} G_3 \right) + \right.

\left. + \left(s_t^\mu p_t^\nu + s_t^\nu p_t^\mu \right) G_6 + i\epsilon^{\mu\nu\rho\sigma} p_{t\rho} s_{t\sigma} G_8 + i\epsilon^{\mu\nu\rho\sigma} q_{\rho} s_{t\sigma} G_9 , \right] \] (4)

There are three unpolarized structure functions \( H_{1,2,3} \) and five polarized structure functions from the set \( G_{1,2,3,6,8,9} \). In general the invariant structure functions are functions of \( q_0 \) and \( q^2 \). In the narrow resonance approximation for the \( W^+–\)boson, which we shall adopt in this talk, one has \( q^2 = m_W^2 \). The aim of the game is to measure the different unpolarized and polarized structure functions (or moments thereof) and to compare them.

\(^1\)In physical expressions the three structure functions \( G_3, G_8 \) and \( G_9 \) contribute only in two pairs of linear combinations [19].
to theoretical predictions. The different structure functions can be separated since they contribute to the rate with different dependencies on the electron energy and, in the case of the polarized structure functions, they can be measured through polar and azimuthal correlations involving the polarization direction of the polarized top quark.

**Unpolarized top quark decays**

In the decay $t \rightarrow X_b + W^+$ the $W^+$ is polarized. The $W^+$ is self–analyzing in the sense that the angular decay distribution of its decay products $W^+ \rightarrow l^+ \nu_l$ can be used to reconstruct the polarization of the $W^+$. We shall analyze the unpolarized decay in cascade–type fashion, i.e. we shall analyze the decay $W^+ \rightarrow l^+ + \nu_l$ in the rest frame of the $W^+$. This brings in the three unpolarized helicity structure functions $H_{T+}, H_{T-}, H_L$ (or for short $H_+, H_-, H_L$) which are linearly related to the three unpolarized invariant structure functions $H_1, H_2, H_3$ via

$$H_+ = H_1 + |\vec{q}| m_t H_3; \quad H_- = H_1 - |\vec{q}| m_t H_3; \quad H_L = m_W^2 H_1 + |\vec{q}|^2 m_t^2 H_2.$$  

The polar angle decay distribution is given by

$$\frac{1}{1-d\cos \theta} = \frac{3}{8}(1 + \cos \theta)^2 H_+ + \frac{3}{8}(1 - \cos \theta)^2 H_- + \frac{3}{4} \sin^2 \theta H_L. \quad (8)$$

where the angle $\theta$ is defined in Fig.2. The $H_+, H_- \text{ and } H_L$ are the normalized transverse–plus, transverse–minus and longitudinal helicity structure functions, resp., such that $H_+ + H_- + H_L = 1$.

From the polar angle dependence in Eq.(8) or from matching $m$–quantum numbers in the $W^+$ rest frame decay (see Fig.2) it is clear that

- $H_+$ : favours forward $l^+$
- $H_-$ : favours backward $l^+$
- $H_L$ : favours forward $l^+$

Translated to the top quark rest frame this implies that $F_+(F_-)$ produce harder (softer) $l^+$'s which can be used to experimentally separate the contributions of the three helicity structure functions.

At the Born term level the SM prediction is ($m_t = 175$ GeV, $m_b = 0$)

$$H_{+\text{(Born)}} = 0 \quad \text{(forbidden)} \quad (9)$$

$$H_{-\text{(Born)}} = \frac{1}{1+2y^2} = 0.297$$

$$H_{L\text{(Born)}} = \frac{2y^2}{1+2y^2} = 0.703,$$

where $y = m_W/m_t$. At the Born term level, with $m_b = 0$, $H_+$ is not populated because of angular momentum conservation in the two–body decay process $t \rightarrow b + W^+$ where for
Figure 2: Definition of the polar angle $\theta$ in the rest frame decay of $W^+ \to l^+ + \nu_l$. The two lines “//” indicate a boost to the rest system of the $W^+$. The arrows next to the lepton lines give the helicities of the leptons.

$m_b = 0$ the bottom quark has 100% negative helicity and the $b$–quark and the $W^+$ are in a back–to–back configuration.

The present experimental results on $\mathcal{H}_+$ are consistent with zero within large error bars. For example, using 230 pb$^{-1}$ D0 quotes a value of $\mathcal{H}_+ = 0.00 \pm 0.13 \text{ (stat)} \pm 0.07 \text{ (syst)}$ \cite{20}. CDF finds $\mathcal{H}_+ = 0.00^{+0.22}_{-0.34} \text{ (stat + syst)}$ or $\mathcal{H}_+ < 0.27$ at the 95% confidence level \cite{21}. Using the same data sample of 200 pb$^{-1}$ CDF quotes a value of $\mathcal{H}_0 = 0.74^{+0.22}_{-0.34}$ for the longitudinal helicity of the $W^+$–boson, also compatible with the SM prediction.

The vanishing of $\mathcal{H}_+$ is no longer true for additional gluon or photon emission, or when one takes into account bottom mass effects. When all of these are taken into account one has \cite{22, 23}

$$\mathcal{H}_+ = 0.00102 \text{(QCD)} + 0.00008 \text{(EW)} + 0.00039 (m_b \neq 0),$$

where the numbers give the $O(\alpha_s)$ QCD corrections, the $O(\alpha)$ electroweak corrections and $m_b \neq 0$ corrections ($m_b = 4.8$ GeV). Numerically the correction to $\mathcal{H}_+$ occurs only at the pro mille level. It is safe to say that, if top quark decays reveal a violation of the SM $(V-A)$ current structure that exceeds the 1% level, the violations must have a non-SM origin.

The results for the corresponding corrections to $\mathcal{H}_-$ and $\mathcal{H}_L$ are listed in terms of rates normalized to the total Born term rate, i.e. $\hat{\Gamma}_i = \Gamma_i / \Gamma(Born)$. The normalized partial Born term rates $\hat{\Gamma}_i(Born)$ are factored out. Corrections coming from NLO QCD, from the NLO electro-weak corrections (EW), from the $W^+$ finite width correction (BW) and from $m_b \neq 0$ effects are listed separately. One has

$$\hat{\Gamma}_- = 0.297 \left[1 - 0.0656 \text{(QCD)} + 0.0206 \text{(EW)} - 0.0197 \text{(BW)} - 0.00172 (m_b \neq 0)\right]$$

$$\hat{\Gamma}_L = 0.703 \left[1 - 0.0951 \text{(QCD)} + 0.0132 \text{(EW)} - 0.0138 \text{(BW)} - 0.00357 (m_b \neq 0)\right]$$

Written in terms of the normalized $\mathcal{H}_i$ this translates into a $+2.4\%$ upward shift from $\mathcal{H}_-(Born) = 0.297$ and a $-1.2\%$ downward shift from $\mathcal{H}_L(Born) = 0.703$. Judging from the fact that $\mathcal{H}_L$ and $\mathcal{H}_-$ will eventually be measured to better than 1% it is quite clear that one has to take radiative corrections into account when comparing experiment with theory.
In Fig. 3 we show the top mass dependence of the ratio $H_L = \Gamma_L/\Gamma$. The horizontal displacement of the Born term curve and the corrected curve is $\approx 3.5$ GeV. One would thus make the corresponding mistake in a top mass determination from the measurement of $H_L$ if the Born term curve was used instead of the corrected curve. If one takes $m_t = 175$ GeV as central value a 1% relative error on $H_L$ would allow one to determine the top quark mass with an error of $\approx 3$ GeV.

4 Polarized top quark decays

Contrary to the analysis of unpolarized top quark decays described in the last subsection polarized top quark decay will be analyzed altogether in the rest frame of the decaying top quark. This is the natural choice for an experimental analysis. Choosing a particular two–particle rest subsystem is only of advantage if that particular subsystem is resonance dominated as was discussed in the unpolarized decay case.

The general angular decay distribution of the rest frame decay of a polarized top quark decaying into a jet $X_b$ and a lepton $l^+$ and a neutrino is given by [24]

$$\frac{d\Gamma}{dx_l d\hat{q}_0 d\cos \theta_P d\phi} = \frac{1}{4\pi} \left( \frac{d\Gamma_A}{dx_l d\hat{q}_0} + P \left( \frac{d\Gamma_B}{dx_l d\hat{q}_0} \cos \theta_P + \frac{d\Gamma_C}{dx_l d\hat{q}_0} \sin \theta_P \cos \phi \right) \right)$$

where the polar and azimuthal angles $\theta_P$ and $\phi$ describe the orientation of the polarization of the top quark relative to the decay plane formed by the decay products of the top quark. The scaled energy and the scaled mass of the $W^+$ are denoted by $\hat{q}_0 = q_0/m_t$ and $y = m_W/m_t$. As usual we define a scaled lepton energy through $x_l = 2E_l/m_t$. $P$ is the magnitude of the top quark polarization. $\Gamma_A$ stands for the unpolarized rate, and $\Gamma_B$ and $\Gamma_C$ stand for the polar and azimuthal correlation rates. In [25] we have considered three different helicity systems to analyse the polar and azimuthal correlations in the rest frame.
It is important to realize that correlation measurements in each of the helicity frames constitute independent measurements of the invariant polarized structure functions. To illustrate this point let us consider the contribution of the invariant polarized structure function $G_1$ to the polar and azimuthal correlations in the above three helicity systems. The decay rate is proportional to $L^{\mu\nu} H_{\mu\nu}$. One then obtains

$$L^{\mu\nu} H_{\mu\nu}(G_1) = m_t q^2 G_1 \left\{ \begin{array}{l}
\frac{x_l \hat{q}_0 - y^2}{x_l} \cos \theta_{P1} + \frac{y}{x_l} \sqrt{x_l(2 \hat{q}_0 - x_l) - y^2 \sin \theta_{P1} \cos \phi} \\
\sqrt{\hat{q}_0^2 - y^2 \cos \theta_{P2}} \\
\frac{\hat{q}_0 - y^2}{2 \hat{q}_0 - x_l} \cos \theta_{P3} + \frac{y}{2 \hat{q}_0 - x_l} \sqrt{x_l(2 \hat{q}_0 - x_l) - y^2 \sin \theta_{P3} \cos \phi} 
\end{array} \right\} ,$$

(12)
where the three contributions in the curly bracket refer to the polar and azimuthal correlations in the three helicity coordinate systems with the z–axes along (1) the lepton \( l^+ \), (2) the \( W^+ \)-boson and (3) the neutrino \( \nu_l \). From Eq.(12) it is clear that \( G_1 \) contributes quite differently to the correlation functions in the three reference systems.

In [25] we have calculated the Born term and NLO QCD contributions to the polar and azimuthal correlation functions \( d\Gamma_B \) and \( d\Gamma_C \) in the three different helicity systems. We were able to obtain closed form expressions for the totally integrated angular decay distributions. The results are too long to be listed here but can be found in [25]. We mention that we find agreement with [27, 28, 29, 30] for the unpolarized case \( d\Gamma_A \) and the polar correlation function \( d\Gamma_B \) in systems 1 and 3. In numerical form one has

- **z–axis along lepton** (system (1a))

\[
\frac{d\Gamma^{\text{NLO}}}{d \cos \theta_P d\phi} = \frac{\Gamma_A^{(0)}}{4\pi} \left[ (1 - 8.54\%) + (1 - 8.72\%) P \cos \theta_P - 0.24\% P \sin \theta_P \cos \phi \right]
\]  

- **z–axis along \( W^+ \)-boson** (system (2'))

\[
\frac{d\Gamma^{\text{NLO}}}{d \cos \theta_P d\phi} = \frac{\Gamma_A^{(0)}}{4\pi} \left[ (1 - 8.54\%) + (0.406 - 11.62\%) P \cos \theta_P - (0.760 - 8.20\%) P \sin \theta_P \cos \phi \right]
\]  

- **z–axis along neutrino** (system (3a))

\[
\frac{d\Gamma^{\text{NLO}}}{d \cos \theta_P d\phi} = \frac{\Gamma_A^{(0)}}{4\pi} \left[ (1 - 8.54\%) - (0.318 - 1.02\%) P \cos \theta_P - (0.919 - 8.61\%) P \sin \theta_P \cos \phi \right]
\]

In all the three expressions we have factored out the Born term rate \( \Gamma_A^{(0)} \). The first number in the round brackets stands for the LO Born term rate whereas the second number gives the percentage change due to the NLO QCD corrections.

Let me first discuss the LO correlation functions. I shall refer to \( \Gamma_B/\Gamma_A \) and \( \Gamma_C/\Gamma_A \) as the polar and azimuthal analyzing power, respectively. In system (1a) \( (l^+ \text{ along } z) \) the polar analyzing power is 100\% which necessarily implies that the azimuthal analyzing power is zero in this system. In fact, the vanishing of \( \Gamma_C \) in system (1a) can be seen to directly follow from the left–chiral \((V - A)\) structure of the SM quark and lepton currents [26]. The polar analyzing power in the systems (2') and (3a) is less than 100\% with +41\% and -32\%, respectively. As mentioned before the LO azimuthal analyzing power in system (1a) is zero. In system (2') and (3a) the azimuthal analyzing power is reasonably large with -76\% and -92\%, respectively.

Except for the polar correlation in system (3a) all NLO corrections go in the same direction. They reduce the LO results by approximately 10\%. This implies that the polar and azimuthal analyzing powers are not changed very much from their Born term values through radiative correction. An exception is system (3a) where the polar analyzing power is changed from -31.8\% to -34.4\%. This amounts to a 8.2\% change in analyzing power through radiative corrections which is surprisingly large.
5 Summary and conclusions

In this talk I have covered three selected topics in top quark physics. The first topic concerned the NNLO calculation of hadronic top quark pair production where the loop–by–loop part is now being completed. The other three missing parts of the NNLO calculation (two–loop, one-loop gluon emission, two–gluon emission) are more difficult and will very likely take another five to ten years to complete. Such a large–size calculation will require a dedicated international effort of the theoretical community which will have to be coordinated by one of the big international centers of particle physics. In the second and third topic I discussed NLO QCD predictions for unpolarized and polarized top quark decays which should be amenable to experimental tests in the next coming few years.

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