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Helicity fractions of W bosons from top quark decays at next-to-next-to-leading order in QCD

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Decay rates of unpolarized top quarks into longitudinally and transversally polarized W bosons are calculated to second order in the strong coupling constant α_s. Including the finite bottom quark mass and electroweak effects, the standard model predictions for the W-boson helicity fractions are $\mathcal{F}_L = 0.687(5)$, $\mathcal{F}_+ = 0.0017(1)$, and $\mathcal{F}_- = 0.311(5)$.

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There has been a continuing interest in the measurement of the helicity fractions of the W boson from top quark decays from the CDF Collaboration [1–7] and from the D0 Collaboration [8–11] at the Tevatron at Fermilab.

In the standard model (SM) the top quark decays predominantly into a W^+ boson and a bottom quark. Interesting observables, independent of the production rate that is difficult to predict precisely for a hadron collider, are the fractions of the three possible W helicities: $\mathcal{F}_L$ (longitudinal), $\mathcal{F}_+$ (transverse plus) and $\mathcal{F}_-$ (transverse minus). In the leading order (LO) in the strong coupling constant α_s (that is, without any gluon corrections), and in the limit of a massless bottom quark one has [12]

$$\mathcal{F}_L{:}\mathcal{F}_+{:}\mathcal{F}_- = \frac{1}{1+2x^2}{:}0{:}\frac{2x^2}{1+2x^2}, \tag{1}$$

with $\mathcal{F}_L + \mathcal{F}_+ + \mathcal{F}_- = 1$ and $x \equiv m_W/m_t$. Using $m_W = 80.401(43)$ GeV [13] and $m_t = 172.8(1.3)$ GeV [14] we get $x^2 = m_W^2/m_t^2 = 0.216(3)$ and $\mathcal{F}_L{:}\mathcal{F}_+{:}\mathcal{F}_- \simeq 0.7{:}0{:}0.3$.

The leading order decay $t \to bW$ is a two-body process. With the $V - A$ interaction, a massless b quark is left handed; thus the W can only be left handed or longitudinal due to angular momentum conservation. One therefore has $\mathcal{F}_+ = 0$ provided no gluons are emitted.

The above LO predictions are only marginally changed by the bottom mass. For a pole mass of $m_b = 4.8$ GeV one finds that the total rate Γ decreases by about a quarter per cent compared to the massless b limit. The helicity fraction $\mathcal{F}_L$ slightly decreases while $\mathcal{F}_-$ increases, by about one per mil. The leakage into the transverse-plus fraction $\mathcal{F}_+$ is less than half per mil. Radiative corrections are a more important source of the transverse-plus rate. However, as we shall see, when NLO and NNLO gluon radiation is included, $\mathcal{F}_+$ still does not exceed two per mil. Since only hard gluon emission can influence the helicity fractions, this smallness is a reliable prediction of the standard model.

For this reason, the transverse-plus fraction $\mathcal{F}_+$ is a sensitive probe of new physics effects such as a right chiral admixture to the SM current. The left and right chiral contributions do not interfere for $m_b = 0$ leading to a quadratic dependence on the admixture parameter. The contribution of the right chiral contribution can be obtained from Eq. (1) by exchanging $\mathcal{F}_+ \leftrightarrow \mathcal{F}_-$, whereas $\mathcal{F}_L$ remains unchanged. We mention that there are some indirect model dependent constraints on a possible right chiral admixture to the SM current from measurements of $b \to s + \gamma$ decays [15–18].

Let us summarize the theoretical prediction for the helicity fractions. In addition to the $\mathcal{O}(\alpha_s^2)$ effects computed in this paper, we include the lower-order contribution [Born + $\mathcal{O}(\alpha_s)$] [19–21], the leading electroweak corrections [22], and, for $\mathcal{F}_+$, the m_b effect. The errors resulting from uncertainties in $m_{t,b,W}$ and α_s and an estimate of the higher-order effects are added in quadrature. We find

$$\mathcal{F}_L = 0.687(5), \qquad \mathcal{F}_+ = 0.0017(1), \qquad \mathcal{F}_- = 0.311(5). \tag{2}$$

The relative errors for $\mathcal{F}_L$ and $\mathcal{F}_-$ are small [of $\mathcal{O}(1\%)$] and, for the largest part, result from the experimental error on the top mass. The error for $\mathcal{F}_+$ arises from uncertainty in α_s and, to a lesser degree, in m_b. Its absolute value is small but the relative error is large due to the fact that $\mathcal{F}_+$ vanishes at LO for $m_b = 0$.

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Various methods have been used by the CDF and D0 Collaborations to experimentally extract the helicity fractions from the top quark decay data (see the recent review [23]). While previous analyses have performed two fits keeping one of the helicity fractions at its SM value, more recent analyses measure the fractions $\mathcal{F}_L$ and $\mathcal{F}_+$ simultaneously in a two-dimensional fit [7,11]. Using such a model independent analysis the CDF Collaboration quotes values of $\mathcal{F}_L = 0.88 \pm 0.11(\mathrm{stat}) \pm 0.06(\mathrm{syst})$ and $\mathcal{F}_+ = -0.15 \pm 0.07(\mathrm{stat}) \pm 0.06(\mathrm{syst})$ [7]. In a similar analysis the D0 Collaboration obtains $\mathcal{F}_L = 0.425 \pm 0.166(\mathrm{stat}) \pm 0.102(\mathrm{syst})$ and $\mathcal{F}_+ = 0.119 \pm 0.090(\mathrm{stat}) \pm 0.053(\mathrm{syst})$ [11]. Both measurements are consistent with the SM predictions.

The experimental errors on the helicity fraction measurements are still rather large but will be much reduced when larger data samples become available in the future from the Tevatron and from the LHC. Optimistically the measurement errors can eventually be reduced to below 1%. For example, an early Monte Carlo (MC) study quotes measurement uncertainties of $\Delta\mathcal{F}_L = 0.007$ and $\Delta\mathcal{F}_+ = 0.003$ for an integrated luminosity of $100\ \mathrm{fb}^{-1}$ at Tevatron II energies [24]. The corresponding event rates can easily be reached at the LHC within one year. A more recent MC study based on $10\ \mathrm{fb}^{-1}$ at the LHC quotes measurement uncertainties of $\Delta\mathcal{F}_L = 0.019$, $\Delta\mathcal{F}_- = 0.018$, and $\Delta\mathcal{F}_+ = 0.0021$ [17].

The improvements in the accuracy of the experimental measurements have to be matched by corresponding advances in the theoretical sector. The NLO $\mathcal{O}(\alpha_s)$ corrections to the helicity fractions were calculated in [19–21]. They lower $\mathcal{F}_L$ and increase $\mathcal{F}_-$ by about 1 and 2%, respectively, relative to their LO values. At NLO there is now a small contribution to the transverse-plus fraction $\mathcal{F}_+$ of 0.001. The corresponding NLO electroweak and finite width corrections were determined in [22]. They are smaller than the strong corrections and tend to cancel each other for both $\mathcal{F}_L$ and $\mathcal{F}_-$.

It is desirable to improve the accuracy of the theoretical predictions and to check the convergence of the perturbative series by computing the helicity fractions at NNLO. A first step in this direction was taken in [25,26],[1] where the NNLO corrections to the total rate were found, exploiting the smallness of $x = m_W/m_t$. A series in powers and logarithms of x was obtained and found to converge rapidly, so that its first few terms suffice. The aim of this paper is to use similar techniques to calculate the NNLO strong corrections to the three helicity fractions.

We first determine the rate Γ_L of the top decay with longitudinally polarized W, replacing the full sum over W polarizations by a projector described below. The previous knowledge of the total rate is used to calculate the transverse rate $\Gamma_T = \Gamma_+ + \Gamma_-$ from the difference $\Gamma_T = \Gamma - \Gamma_L$. We use another projector to find the difference $\Gamma_+ - \Gamma_-$. Finally, the helicity fractions $\mathcal{F}_i = \Gamma_i/\Gamma$ are determined.

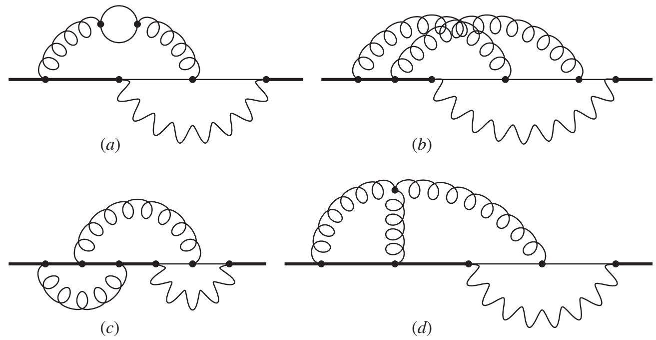

FIG. 1. Sample three-loop diagrams. Thick and thin lines denote top and bottom quarks, respectively. Wavy lines denote W bosons and curly lines denote gluons. In the closed fermion loop all quark flavors have to be considered.

Our calculation follows the approach outlined in Ref. [26]. Using the optical theorem we compute the decay width from the imaginary part of top quark self-energy diagrams,

$$\Gamma = \frac{1}{m_t}\,\mathrm{Im}(\Sigma), \tag{3}$$

where Σ denotes the one-particle irreducible self-energy diagrams. Sample diagrams are shown in Fig. 1. The unitary gauge is used for the W boson so that diagrams with Goldstone bosons are not needed. However, the R_ξ gauge is used for the gluons with an arbitrary gauge parameter. The gauge-parameter dependence cancels in the final result.

Since we set the mass of the bottom quark to zero, the integrals contain two scales, m_t and m_W. To reduce these integrals to single scale integrals, we use the method of expansion by regions (see, e.g., Ref. [28]). In the present case, there are two regions to be considered. In the so-called hard region, the loop momenta are of the order of m_t, while they are of order m_W in the so-called soft region. The integrals become scaleless and vanish if a gluon momentum is soft. Thus, we are left with two contributions to each integral: one where all momenta are hard and one where only the W-boson momentum is soft. For each contribution we construct appropriate expansions in the corresponding small quantities. The remaining single scale integrals are further reduced to so-called master integrals using Laporta's algorithm [29,30].

Compared to the NNLO calculation of the total width, to get the partial rates with various W polarizations requires replacing the total rate projector[2]

[1] An approximate value for the total rate was found in Ref. [27].

[2] In the unitary gauge, the W-boson propagator reads $i\mathbb{P}^{\mu\nu}/(q^2 - m_W^2)$.

$$\mathbb{P}^{\mu\nu} = -g^{\mu\nu} + \frac{q^\mu q^\nu}{m_W^2}, \tag{4}$$

by the longitudinal projector $\mathbb{P}_L^{\mu\nu}$ or the transverse-plus/minus projectors $\mathbb{P}_\pm^{\mu\nu}$. The longitudinal projector reads [20]

$$\mathbb{P}_L^{\mu\nu} = \frac{(m_W^2 p^\mu - p\cdot q q^\mu)(m_W^2 p^\nu - p\cdot q q^\nu)}{m_W^2 m_t^2 |\vec{q}|^2}, \tag{5}$$

where p is the top quark momentum and $q = (q_0; 0, 0, |\vec{q}|)$ is the momentum of the W boson which propagates in the z direction. The transverse projectors can be obtained with the help of the forward-backward projector [20]

$$\mathbb{P}_F^{\mu\nu} = -\frac{1}{m_t|\vec{q}|} i\varepsilon^{\mu\nu\sigma\rho} p_\sigma q_\rho. \tag{6}$$

One has $\mathbb{P}_\pm^{\mu\nu} = (\mathbb{P}^{\mu\nu} - \mathbb{P}_L^{\mu\nu} \pm \mathbb{P}_F^{\mu\nu})/2$.

Partial helicity rates involve two technical challenges absent in the total rate calculation. First, there is an additional propagatorlike structure $1/|\vec{q}|^n$, $n \in \{1, 2\}$ in Eqs. (5) and (6), and second, we have to deal with the presence of γ_5-odd traces in dimensional regularization. Our approach to both issues is outlined below.

In the hard region, we express $|\vec{q}|^2$ through the propagator factor $N = (p+q)^2 - m_t^2 = 2pq + q^2$ as follows[3]:

$$\begin{aligned} |\vec{q}|^2 = q_0^2 - m_W^2 &= \frac{(2p\cdot q)^2}{4m_t^2} - m_W^2 \\ &= \frac{1}{4m_t^2}[N^2 - 2m_W^2 N + m_W^4 - 4m_t^2 m_W^2]. \end{aligned} \tag{7}$$

In Eq. (7) we use the fact that we are only interested in the imaginary part and that $q^2 = m_W^2$ on the cut. Now we can construct the desired expansions in m_W/m_t as

$$\begin{aligned} \frac{1}{|\vec{q}|} &= \frac{2m_t}{N}\sum_{i=0}^{\infty}\binom{2i}{i}\left(\frac{2m_W^2 N - m_W^4 + 4m_t^2 m_W^2}{4N^2}\right)^i, \\ \frac{1}{|\vec{q}|^2} &= \frac{4m_t^2}{N^2}\sum_{i=0}^{\infty}\left(\frac{2m_W^2 N - m_W^4 + 4m_t^2 m_W^2}{N^2}\right)^i, \end{aligned} \tag{8}$$

which we truncate at some order. Thus, the additional propagatorlike structure from the projector is transformed into a scalar on-shell propagator with momentum $p + q$ and mass m_t, raised to arbitrary, integer powers. For the calculation of the polarized decays we need, next to the master integrals of Refs. [26,31], twelve additional three-loop master integrals.

In the soft region, we cannot perform an expansion of $|\vec{q}|$, since $|\vec{q}|^2 = q_0^2 - m_W^2$ and q_0 is of order m_W in the soft region. However, in this region the W-boson loop factorizes. Therefore, we only have to replace the usual one-loop tadpole integrals with integrals of the type

$$\int \frac{d^d q}{(q^2 - m_W^2)(q_0^2 - m_W^2)^n}, \tag{9}$$

with $n \in \{1/2, 1\}$. $d = 4 - 2\epsilon$ is the number of dimensions. Integrals of this type can be easily evaluated by performing the integrations over the timelike and spacelike momentum components separately.

For traces with an odd power of γ_5, we use the prescription of Ref. [32] and replace

$$\gamma_\mu\gamma_5 \to \frac{i}{3!}\varepsilon_{\mu\alpha\beta\delta}\gamma^\alpha\gamma^\beta\gamma^\delta. \tag{10}$$

The ε tensor is stripped off and absorbed into the projector. As a consequence the renormalization constant of the axial-vector current at the requisite order becomes

$$Z_A = 1 + \left(\frac{\alpha_s}{\pi}\right)^2\left(\frac{11}{24}C_F C_A - \frac{1}{6}C_F T_F n_f\right)\frac{1}{\epsilon}, \tag{11}$$

where $C_F = (N_c^2 - 1)/(2N_c)$ and $C_A = N_c$ are the Casimir operators of the fundamental and adjoint representation of SU(N_c), respectively. For QCD we have $N_c = 3$ and $T_F = 1/2$. n_f denotes the number of quark flavors. Additionally, we have to include the finite renormalization constant

$$\begin{aligned} Z_5 = 1 - \frac{\alpha_s}{\pi}C_F + \left(\frac{\alpha_s}{\pi}\right)^2 \\ \times\left(\frac{11}{8}C_F^2 - \frac{107}{144}C_F C_A + \frac{1}{36}C_F T_F n_f\right) \end{aligned} \tag{12}$$

to restore the anticommutativity of γ_5. Both renormalization constants were determined at the three-loop level in Ref. [33].

A NLO check of the new methods used in this paper is afforded by comparing with the expanded form of the known NLO closed form results given in [19–21]. We found agreement up to $\mathcal{O}(x^{16})$.

We present our results in terms of the reduced helicity rates $\hat{\Gamma}_i$ where

$$\Gamma_i = \frac{G_F m_t^3 |V_{tb}|^2}{8\sqrt{2}\pi}\hat{\Gamma}_i \tag{13}$$

with $i \in \{L, +, -\}$. G_F is Fermi's constant and V_{tb} is the element of the Cabibbo-Kobayashi-Maskawa matrix which governs transitions between bottom and top quarks.

The analytical results of our calculation are too long to be presented here. Instead we present their numerical values. In Table I we show successive terms in the power series expansion [in terms of $[x^n] := (x^n, x^n \ln x)$] of the NNLO correction to the reduced helicity rates up to terms of order $\mathcal{O}([x^{10}])$. Noteworthy is the fact that there are also odd powers in the expansion in $[x]$. These terms appear in the expansion of the parity-odd helicity structure function

[3] $p = (m_t; 0, 0, 0)$ in the rest frame of the top quark such that $p\cdot q = m_t q_0$.

TABLE I. Numerical values for $\mathcal{O}([x^n] := (x^n, x^n \ln x))$ terms in the x expansion of the NNLO corrections to the reduced partial and total helicity rates $\hat{\Gamma}_i$.

	$\hat{\Gamma}_L$	$\hat{\Gamma}_+$	$\hat{\Gamma}_-$	$\hat{\Gamma}$
$[x^0]$	-1.958×10^{-2}	0	0	-1.958×10^{-2}
$[x^2]$	4.737×10^{-3}	3.860×10^{-4}	-3.861×10^{-3}	1.262×10^{-3}
$[x^4]$	6.710×10^{-4}	1.351×10^{-4}	-9.917×10^{-4}	-1.856×10^{-4}
$[x^5]$	0	-5.339×10^{-4}	5.339×10^{-4}	0
$[x^6]$	-1.467×10^{-4}	1.186×10^{-4}	4.878×10^{-4}	4.597×10^{-4}
$[x^7]$	0	7.696×10^{-5}	-7.696×10^{-5}	0
$[x^8]$	-1.702×10^{-5}	-2.333×10^{-5}	-1.723×10^{-5}	-5.758×10^{-5}
$[x^9]$	0	3.408×10^{-6}	-3.408×10^{-6}	0
$[x^{10}]$	-1.274×10^{-6}	-1.844×10^{-6}	-2.176×10^{-6}	-5.294×10^{-6}
Σ	-1.434×10^{-2}	1.610×10^{-4}	-3.931×10^{-3}	-1.811×10^{-2}

Γ_F[4] and thereby in $\Gamma_\pm$ (see Table I). The odd powers of x stem solely from the soft region of Γ_F. The leading x^0 contributions of $\hat{\Gamma}_L$ and $\hat{\Gamma}$ are equal to each other. This is a consequence of the Goldstone boson equivalence theorem. Between the $[x^4]$ and $[x^6]$ terms the power series expansion is somewhat erratic for $\hat{\Gamma}_+$, $\hat{\Gamma}_-$, and $\hat{\Gamma}$. However, expanding up to $\mathcal{O}([x^{10}])$ Table I shows that one has sufficient numerical stability and precision for all three helicity rates and their sum. The contribution of the $\mathcal{O}([x^{10}])$ term amounts to about 0.01%, 1%, 0.06%, and 0.03% of the total for $\hat{\Gamma}_L$, $\hat{\Gamma}_+$, $\hat{\Gamma}_-$, and $\hat{\Gamma}$, respectively. The convergence is slowest for $\hat{\Gamma}_+$. But then $\hat{\Gamma}_+$ is numerically very small.

In order to present our numerical results on the helicity fractions we define helicity fractions up to $\mathcal{O}(n)$ by writing ($n = 0, 1, 2$ denote the contributions up to LO, NLO, and NNLO, respectively)

$$\mathcal{F}_i^{(n)} = \frac{\sum_{j=0}^n \Gamma_i^{(j)}}{\sum_{j=0}^n \Gamma^{(j)}}, \tag{14}$$

where $i = L, +, -$. We further define the increments $\Delta\mathcal{F}_i^{(n)} = \mathcal{F}_i^{(n)} - \mathcal{F}_i^{(n-1)}$ and the relative increments $\delta\mathcal{F}_i^{(n)} = \Delta\mathcal{F}_i^{(n)}/\mathcal{F}_i^{(0)}$. We present our numerical results in the form $\mathcal{F}_i = \mathcal{F}_i^{(0)} + \Delta\mathcal{F}_i^{(1)} + \Delta\mathcal{F}_i^{(2)}$, and also, if $\mathcal{F}_i^{(0)} \neq 0$, as $\mathcal{F}_i = \mathcal{F}_i^{(0)}(1 + \delta\mathcal{F}_i^{(1)} + \delta\mathcal{F}_i^{(2)})$. For our numerical results we use $\alpha_s(m_t) = 0.1073(24)$, which we obtained with the program RUNDEC [34] from the values $\alpha_s(m_Z) = 0.1176(20)$ and $m_Z = 91.1876(21)$ GeV [35]. We find

$$\begin{aligned}
\mathcal{F}_L &= 0.6978 - 0.0075 - 0.0023 \\
&= 0.6978(1 - 0.0108 - 0.0033), \\
\mathcal{F}_+ &= 0 + 0.001\,03 + 0.000\,23, \\
\mathcal{F}_- &= 0.3022 + 0.0065 + 0.0021 \\
&= 0.3022(1 + 0.0215 + 0.0070).
\end{aligned} \tag{15}$$

The results in Eq. (15) contain higher orders in α_s from the expansion of the denominators in Eq. (14). In order to maintain the constraint $\mathcal{F}_L + \mathcal{F}_+ + \mathcal{F}_- = 1$, we prefer the unexpanded definition of helicity fractions (14).

The numbers in Eq. (15) show the good convergence of the perturbative expansion, even though $\Delta\mathcal{F}_i^{(1)}/\mathcal{F}_i^{(0)}$ (for $i = L, -$) is much smaller than $\Delta\mathcal{F}_i^{(2)}/\Delta\mathcal{F}_i^{(1)}$. The NLO corrections to the helicity fractions are already close to the expected future experimental sensitivities and the NNLO corrections increase these by approximately a third. In particular, the NNLO calculation of the helicity fraction $\mathcal{F}_+$ remains at the order of 0.001. Should a measurement reveal a significantly larger value, it would be a clear signal of new physics.

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[4]This follows the pattern in unpolarized and polarized top quark decays where the expansion of the five parity-even structure functions have $n =$ even whereas the expansion of the five parity-odd structure functions have $n =$ even/odd [20].

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