Abstract

The three-loop QCD contributions to the vacuum polarization functions of the Z and W bosons at zero momentum are calculated. The top quark is considered to be massive and the other quarks massless. Using these results, we calculate the correction to the electroweak $\rho$ parameter.

All computations are done in the framework of dimensional regularization as well as regularization by dimensional reduction. We use recurrence relations obtained by the method of integration by parts to reduce all integrals to a small set of master integrals.

A comparison of the two-loop and three-loop QCD corrections to the $\rho$ parameter is performed.
Owing to the apparent discovery of the top quark [1] with a mass of $174 \pm 10$ GeV, the prospects grow to test the Standard Model on an even higher level of precision than it was possible by now. In particular, the precise knowledge of top-mass effects will allow us to obtain better limits on virtual Higgs effects (and thus, indirectly, on the Higgs mass), and possibly, on new physics. For this reason, a great deal of work has been devoted to the study of the top-mass effects in higher-loop radiative corrections of various electroweak parameters. Mostly, in these studies the top mass is assumed to be large compared to all other masses so that the latter can be put equal to zero from the very beginning (see [2]). In [3, 4] also the Higgs mass $m_H$ was kept as an independent parameter, and both limits, $m_H \ll m_t$ and $m_H \gg m_t$, were studied.

In the Standard Model there are two different sources of corrections which become large ($\sim G_\mu m_t^2$) in the limit of a heavy top, owing to the large top-bottom mass ratio: the $Z$ and $W$ self-energies (affecting, in particular, the $\rho$ parameter) [5] and the $Zb\bar{b}$ vertex [6]. Experimentally, these effects are best accessible in $e^+e^- \to f\bar{f}$ near the $Z$ resonance measured at LEP1 and the on-resonance asymmetries measured at LEP1 and the SLAC $e^+e^-$ linear collider SLC.

In the present paper we are concerned with the heavy-top QCD-corrections to the electroweak $\rho$ parameter in three-loop approximation. The $\rho$ parameter is defined as the ratio of the neutral-current to charged-current amplitudes at zero momentum transfer:

$$\rho = \frac{G_{NC}(0)}{G_{CC}(0)} = \frac{1}{1 - \Delta \rho},$$

where the leading fermion contribution to $\Delta \rho$ is contained in the gauge-boson self-energies

$$\Delta \rho = \frac{\Pi_Z(0)}{M_Z^2} - \frac{\Pi_W(0)}{M_W^2}.$$  \hspace{1cm} (2)

In the approximation considered, we write

$$\Delta \rho = 3x_t(1 + \delta^{EW} + \delta^{EW,QCD} + \delta^{QCD}) \simeq 3x_t(1 + \rho^{(2)} x_t)(1 + h\delta^{QCD} + h^2 \delta^{QCD}(3)),$$  \hspace{1cm} (3)

with

$$x_t = \frac{\sqrt{2} G_\mu}{16\pi^2} m_t^2, \hspace{1cm} h = \frac{\alpha_s}{4\pi},$$  \hspace{1cm} (4)

$\alpha_s$ being the QCD coupling constant. We have denoted by $\delta^{EW}$ the pure electroweak, by $\delta^{EW,QCD}$ the mixed electroweak-QCD, and by $\delta^{QCD}$ the pure QCD corrections. The two-loop electroweak correction $\rho^{(2)}$, due to virtual Higgs (ghost) effects, is small $\rho^{(2)}|_{m_H=0} = -0.74$ for $m_H \approx 0$ [7] but reaches a maximum as large as $-11.57$ at $m_H \approx 5.7 m_t$ [3, 4].

The one-loop correction to $\Delta \rho$ was first calculated in [5]. The two-loop QCD correction $\delta^{QCD}_{(2)}$ has been calculated in [8]. It proved to be rather large. If one takes $m_t$ as the top-quark pole mass, then

$$\delta^{QCD}_{(2)} = -\frac{8}{9}(\pi^2 + 3).$$  \hspace{1cm} (5)

Therefore, it is essential to evaluate the next, three-loop correction, in view of the high precision of modern experiments.
To evaluate $\Delta \rho$, the diagonal parts of the self-energies of the $W$ and $Z$ gauge bosons

$$\Pi^{\mu\nu}_\alpha(q) = g_{\mu\nu} \Pi_\alpha(q^2) + g_{\mu\nu} q_\nu \tilde{\Pi}_\alpha(q^2)$$

($\alpha = W, Z$) at $q = 0$ are needed. Since at zero momentum and $m_t \neq 0$ no infrared divergences appear in diagrams with only fermions and gluons, one may put $q = 0$ from the very beginning. Contracting $\Pi^{\mu\nu}_\alpha(0)$ with $g_{\mu\nu}$, we obtain for $\Pi_\alpha(0)$ an expression containing only bubble integrals. At the one- and two-loop level these are quite simple and for arbitrary space-time dimension $d = 4 - 2\varepsilon$ can be written in terms of Euler’s $\Gamma$ function. Here we need only

$$\int \frac{d^dk_1}{\pi^{d/2}} \frac{(m^2)^{\beta-d/2}}{(k_1^2 + m^2)^\beta} = \frac{\Gamma(\beta - d/2)}{\Gamma(\beta)},$$

$$\int \frac{d^dk_1 d^dk_2}{\pi^d} \frac{(m^2)^{\alpha+\beta+\gamma-d}}{(k_1^2 + m^2)^\alpha(k_2^2 + m^2)^\beta ((k_1 - k_2)^2)^\gamma} = \frac{\Gamma(\alpha + \beta + \gamma - d) \Gamma(\frac{d}{2} - \gamma) \Gamma(\beta + \gamma - \frac{d}{2}) \Gamma(\alpha + \gamma - \frac{d}{2})}{\Gamma(\alpha) \Gamma(\beta) \Gamma(\gamma) \Gamma(d/2)},$$

$$\int \frac{d^dk_1 d^dk_2}{\pi^d} \frac{(m^2)^{\alpha+\beta+\gamma-d}}{(k_1^2 + m^2)((k_1 - k_2)^2)^\beta(k_2^2 + m^2)^\gamma} = \frac{\Gamma(\alpha + \beta + \gamma - d) \Gamma(\alpha + \beta - d/2) \Gamma(d/2 - \alpha) \Gamma(d/2 - \beta)}{\Gamma(\alpha) \Gamma(\beta) \Gamma(\gamma) \Gamma(d/2)}. \quad (7)$$

At the three-loop level 22 diagrams of the $Z$-boson self-energy and 29 diagrams of the $W$-boson self-energy contribute to $\delta^{QCD}$. The integrals that appear here are much more complicated than at the one- and two-loop level.

The rather complicated task of computing massive three-loop Feynman diagrams is accomplished by applying the method of recurrence relations [1, 11]. This method allows us to relate various scalar Feynman integrals of the same prototype which differ by powers of their scalar propagators. As a result, by means of plain algebra, any diagram is reduced to a limited number of so-called master integrals. They need to be evaluated once and for all, and can then be used in any renormalizable quantum field theory. Some of the integrals that we need for the present three-loop calculation were considered in [11]. Here, however, more types of integrals are required. In addition to the master integrals evaluated in [11], two more nontrivial master integrals are encountered:

$$\int \int \int \frac{d^dk_1 d^dk_2 d^dk_3}{\pi^{d/2} \Gamma(1 + \varepsilon)^3} \frac{(m^2)^{4-3d/2}}{k_1^2[(k_1 - k_2)^2 + m^2][(k_2 - k_3)^2 + m^2][k_3^2 + m^2]}$$

$$= \frac{1}{\varepsilon^3} + \frac{15}{4\varepsilon^2} + \frac{65}{8\varepsilon} + \frac{135}{16} + \frac{81}{4} S_2 + O(\varepsilon), \quad (8)$$

$$\int \int \int \frac{d^dk_1 d^dk_2 d^dk_3}{\pi^{d/2} \Gamma(1 + \varepsilon)^3} \frac{(m^2)^{6-3d/2}}{k_1^2(k_1 - k_2)^2k_3^2[(k_1 - k_2)^2 + m^2][k_2^2 + m^2][(k_2 - k_3)^2 + m^2]}$$

$$= 2\zeta(3) \frac{1}{\varepsilon} + D_3 + O(\varepsilon) \quad (9)$$
with
\[ S_2 = \frac{4}{9\sqrt{3}} \text{Cl}_2 \left( \frac{\pi}{3} \right) = 0.260434137632162098955729. \]  \hspace{1cm} (10)

We have not found a representation of \( D_3 \) in terms of known transcendental numbers, though we do not exclude its existence. By means of the numerical method for the evaluation of Feynman diagrams proposed in [12], \( D_3 \) can be calculated quite accurately. Here we give the first 22 digits, which is more than enough for a precise evaluation of \( \rho \):
\[ D_3 = -3.027009493987652019786. \]  \hspace{1cm} (11)

Calculations were mostly done using FORM 1.1 [11]. All the diagrams were computed in the covariant gauge with an arbitrary gauge parameter. Performing charge and mass renormalization in the \( \overline{MS} \) scheme, we got for the \( W \)-boson propagator the following expression:
\[
\Pi^{(3)}_W(0) = 12x_tM^2_W \left\{ \left( -\frac{1}{2\varepsilon} - \frac{1}{4} - \frac{1}{2}\hat{l} \right) + C_F \left( \frac{3}{2\varepsilon^2} - \frac{5}{4\varepsilon} - \frac{13}{8} + \zeta(2) - \frac{1}{2}\hat{l} - \frac{3}{2}\hat{l}^2 \right) \right\} \hspace{1cm} (12)
\]
Here \( n_f \) is the total number of quarks, \( \hat{l} = \ln(\mu^2/\hat{m}_t^2(\mu)) \), \( \hat{m}_t(\mu) \) is the renormalized mass in the \( \overline{MS} \) scheme, and the constant
\[ B_4 = 16 \text{Li}_4 \left( \frac{1}{2} \right) + \frac{2}{3} \ln^4 2 - \frac{2}{3} \pi^2 \ln^2 2 - \frac{13}{180} \pi^4 = -1.762800087073770086 \]  \hspace{1cm} (13)
was defined in [11]. The result for the \( Z \)-boson propagator is
\[
\Pi^{(3)}_Z(0) = 12x_tM^2_Z \left\{ \left( -\frac{1}{2\varepsilon} - \frac{1}{2}\hat{l} \right) + C_F \left( \frac{3}{2\varepsilon^2} - \frac{5}{4\varepsilon} - \frac{1}{8} + \frac{1}{2}\hat{l} - \frac{3}{2}\hat{l}^2 \right) \right\} \hspace{1cm} (12)
\]
\[
(\begin{array}{c}
+\left(\frac{101}{8} - 18\zeta(3)\right)i + \frac{39}{4}i^2 - 3i^3 \\
+C_A C_F \left(-\frac{11}{6\varepsilon^3} + \frac{83}{12\varepsilon^2} - \frac{77}{12\varepsilon} + \frac{3}{\varepsilon}\zeta(3) + 3 + \frac{28}{3}\zeta(3) - \frac{27}{2}\zeta(4) + 3B_4 \right) \\
+ \left(-\frac{85}{24} + 9\zeta(3)\right)i - \frac{43}{6}i^2 - \frac{11}{6}i^3 \\
+C_F n_f \left(\frac{1}{3\varepsilon^3} - \frac{5}{6\varepsilon^2} - \frac{1}{12} + \frac{8}{3}\zeta(3) + \frac{5}{12}i + \frac{2}{3}i^2 + \frac{1}{3}i^3 \right) \\
+C_F \left(-2 - 12\zeta(3)\right) \right) h^2.
\end{array}\} \tag{14}
\]

In the sum of the bare diagrams contributing to \(\Pi_\alpha(0)\) \((\alpha = W, Z)\), as well as in the corresponding counterterms separately, the gauge parameter cancels, which is a partial check of our result.

Substitution of (12) and (14) into (2) gives the ultraviolet-finite expression

\[
\delta^{QCD}_{(3),\overline{MS}} = C_F^2 \left[ -\frac{1591}{36} - \frac{518}{9}\zeta(2) + \frac{1084}{3}\zeta(3) + 4\zeta(4) - 1053S_2 + 4D_3 + 8B_4 \\
+\left(3 - 24\zeta(2)\right)i + 18i^2 \right]
\]

\[
+C_A C_F \left[\frac{1013}{12} - \frac{146}{3}\zeta(2) - \frac{452}{3}\zeta(3) + 30\zeta(4) + \frac{1053}{2}S_2 \\
-2D_3 - 4B_4 + \left(\frac{163}{3} - \frac{44}{3}\zeta(2)\right)i + 11i^2 \right]
\]

\[
+C_F n_f \left[ -\frac{25}{2} + \frac{28}{3}\zeta(2) - \frac{16}{3}\zeta(3) + \left(-\frac{22}{3} + \frac{8}{3}\zeta(2)\right)i - 2i^2 \right]
\]

\[
+C_F \left[\frac{152}{3} - \frac{28}{3}\zeta(2) - \frac{512}{3}\zeta(3) + 486S_2 \right]. \tag{15}
\]

If we perform mass renormalization in such a way that the renormalized mass is the pole mass \(m_t\), then we obtain

\[
\delta^{QCD}_{(3)} = C_F^2 \left[ -\frac{238}{9} - \frac{770}{9}\zeta(2) + 96\zeta(2) \ln 2 + \frac{1012}{3}\zeta(3) + 4\zeta(4) \\
-1053S_2 + 4D_3 + 8B_4 \right]
\]

\[
+C_A C_F \left[ -\frac{49}{6} - \frac{98}{3}\zeta(2) - 48\zeta(2) \ln 2 - \frac{416}{3}\zeta(3) + 30\zeta(4) + \frac{1053}{2}S_2 \\
-2D_3 - 4B_4 + \left(-\frac{22}{3} - \frac{44}{3}\zeta(2)\right)i \right]
\]
\[ +C_F n_f \left[ \frac{2}{3} + \frac{52}{3} \zeta(2) - \frac{16}{3} \zeta(3) + \left( \frac{4}{3} + \frac{8}{3} \zeta(2) \right) l \right] + C_F \left[ \frac{188}{3} - \frac{100}{3} \zeta(2) - \frac{512}{3} \zeta(3) + 486.5 \right]. \]  

(16)

In the above formula \( l = \ln(\mu^2/m_t^2) \). Ultraviolet finiteness of (15) and (16) is an additional check of our result. Expression (16) can also be obtained from our result in the \( \overline{MS} \) scheme by using the relation between \( \hat{m}_t \) in the \( \overline{MS} \) scheme and the pole mass \( m_t \) \[13\].

In the \( \overline{MS} \) scheme for QCD [i.e. for the SU(3) gauge group with \( C_A = 3, C_F = 4/3 \)] we get the following expression:

\[ \delta_{QCD}^{\overline{MS}} = \left( 8 - \frac{16}{3} \zeta(2) + 8 l \right) h + \left[ \frac{26459}{81} - \frac{16}{9} B_4 - \frac{25064}{81} \zeta(2) - \frac{5072}{27} \zeta(3) \right] + \frac{1144}{9} \zeta(4) + 882 S_2 - \frac{8}{9} D_3 + n_f \left( -\frac{50}{3} + \frac{112}{9} \zeta(2) - \frac{64}{9} \zeta(3) \right) \]

\[ + \left( \frac{668}{3} - \frac{304}{3} \zeta(2) + \left( -\frac{88}{9} + \frac{32}{9} \zeta(2) \right) n_f \right) \hat{l} + \left( 76 - \frac{8}{3} n_f \right) \hat{l}^2 \right] h^2. \]  

(17)

Substituting numerical values for all the constants and taking \( \mu^2 = \hat{m}_t^2 \) with the minimally subtracted mass we obtain

\[ \delta_{QCD}^{\overline{MS}} = -0.061 511 928 430 2 \alpha_s - (0.221 937 307 314 4 + 0.030 043 860 323 8) \alpha_s^2, \]  

(18)

which at \( n_f = 6 \) turns into

\[ \delta_{QCD}^{\overline{MS}} = -0.061 511 928 430 2 \alpha_s + 0.402 200 469 257 5 \alpha_s^2. \]  

(19)

The smallness of this correction in the \( \overline{MS} \) scheme confirms expectations about higher order effects in electroweak parameters (see e.g. \[14\]).

With the definition of the renormalized mass \( m_t \) as the pole mass of the top quark we obtain:

\[ \delta_{QCD} = -\frac{2}{3} \left( 1 + 2 \zeta(2) \right) \frac{\alpha_s}{\pi} + \left[ \frac{157}{648} - \frac{3313}{162} \zeta(2) - \frac{308}{27} \zeta(3) \right] \]

\[ + \frac{143}{18} \zeta(4) - \frac{4}{3} \zeta(2) \ln 2 + \frac{441}{8} S_2 - \frac{1}{9} B_4 - \frac{1}{18} D_3 \]

\[ - \left( \frac{1}{18} - \frac{13}{9} \zeta(2) + \frac{4}{9} \zeta(3) \right) n_f - \left( \frac{11}{6} - \frac{1}{9} n_f \right) \left( 1 + 2 \zeta(2) \right) l \right] \left( \frac{\alpha_s}{\pi} \right)^2. \]  

(20)

Substituting numerical values for all the constants and putting \( \mu^2 = m_t^2 \) we get

\[ \delta_{QCD} = -0.910 338 291 586 9 \alpha_s - (2.564 571 412 664 2 - 0.180 981 195 767 9) \alpha_s^2, \]  

(21)

and at \( n_f = 6 \) we have

\[ \delta_{QCD} = -0.910 338 291 586 9 \alpha_s - 1.478 684 237 779 7 \alpha_s^2. \]  

(22)
Both two- and three-loop QCD contributions are negative, and their effect is a screening of the bare mass splitting, so that only a reduced “effective” quantity enters the $\rho$ parameter.

Following the method of fastest apparent convergence [15], we can absorb our three-loop correction into a rescaling of $\alpha_s$. With $n_f = 6$, $\delta^{(3)}_{QCD}$ will be zero, if we take $\mu \approx 0.2327 m_t$. When we apply the BLM procedure [16], the $n_f$-dependent term in (20) can be absorbed into the rescaling of $\alpha_s$, if we choose $\mu \approx 0.154 m_t$. The same value was also obtained in [17]. We conclude that the expression for the term proportional to $n_f$, given in [17], agrees with the analytical result for this term given in (20). As an important result of our calculation, we stress the stability of $\delta^{QCD}$: the usual perturbation theory, the FAC and BLM procedures give rather close results for $\delta^{QCD}$. Taking $\alpha_s(m_t) = 0.1055$, we have $\delta^{QCD} = -0.1125$, $-0.1159$ and $-0.1154$ for the perturbation theory, FAC and BLM procedures, respectively.

For our calculations we used the anticommuting $\gamma_5$. To check our result obtained with this prescription at least partially, we calculated $\delta^{QCD}$ again, using the regularization by dimensional reduction [18] which keeps the algebra of $\gamma$ matrices four-dimensional. For propagator-type diagrams in the three-loop approximation the inconsistency of this recipe is not yet revealed. As was expected, the result that we obtained in the regularization by dimensional reduction agrees with $\delta^{QCD}$ in the conventional dimensional regularization after the following recalculation of the coupling constant:

$$h_{RDR} = h \left(1 + \frac{1}{3} C_A h\right),$$  

which just corresponds to a change in the renormalization scheme. The relation can be derived, for example, by equating invariant charges in these two schemes.

Special attention was paid to the evaluation of the diagram with the axial anomaly. Only one such diagram contributes to $\Delta \rho$ (namely, to $\Pi_\rho(0)$) with two triangles and only the top quark running around. We evaluated it in the framework of the regularization by dimensional reduction which keeps the algebra of $\gamma$ matrices four-dimensional. For propagator-type diagrams in the three-loop approximation the inconsistency of this recipe is not yet revealed. As was expected, the result that we obtained in the regularization by dimensional reduction agrees with $\delta^{QCD}$ in the conventional dimensional regularization after the following recalculation of the coupling constant:

$$h_{\rho} = h \left(1 + \frac{1}{3} C_A h\right),$$

(23)

Some observables that are affected by our result are briefly mentioned here. One of them is the mass of the $W$ boson as predicted from $\alpha$, $G_\mu$ and $M_Z$ [21]

$$M_W^2 = \frac{\rho M_Z^2}{2} \left(1 + \sqrt{1 - \frac{4 A_0^2}{\rho M_Z^2} \left(\frac{1}{1 - \Delta \alpha} + \ldots\right)}\right),$$

(24)

where $A_0 = \left(\frac{\pi \alpha}{\sqrt{2} G_\mu}\right)^{1/2} = 37.2802(3)$ and $\Delta \alpha \simeq 0.06$ is the shift of the fine structure constant $\alpha$ due to photon vacuum polarization effects. The ellipsis stands for the non-leading remainder terms. Another physical quantity is the effective weak mixing parameter relevant to Z-resonance physics. It is given by

$$\sin^2 \Theta = 1 - \frac{M_W^2}{\rho M_Z^2} = \frac{1}{2} \left(1 - \frac{4 A_0^2}{\rho M_Z^2} \left(\frac{1}{1 - \Delta \alpha} + \ldots\right)\right),$$

(25)
Numerical values illustrating how these observables are affected by various corrections are given in Table 1. QCD corrections were calculated using (22).

Table 1: Percentage of various two- and three-loop QCD heavy-top corrections at $\alpha_s(m_t) = 0.1055$ and $m_t = 174$ GeV. The value of $\alpha_s$ was obtained by extrapolating $\alpha_s(M_Z=91.1895$ GeV)$=0.118$ to the scale $\mu = m_t$ with the aid of the three-loop $\beta$ function [22] with $n_f = 5$.

| Observable | Two-loop electroweak | Three-loop QCD | Two-loop electroweak |
|------------|----------------------|----------------|---------------------|
| $M_W$      | $-0.065$             | $-0.011$       | $-0.002$           |
| $\sin^2 \Theta$ | $0.130$             | $0.022$       | $0.003$           |

Table 1 demonstrates that the three-loop QCD correction is comparable with the two-loop electroweak correction for sizable Higgs masses (for $m_H/m_t=1.5$ the former amounts to more than 60% of the latter). Thus, we conclude that it makes sense to evaluate subleading electroweak two-loop corrections to $\Delta \rho$ (or another observable) only if the three-loop QCD corrections are taken into account as well.

Acknowledgments. L.Avdeev, S.Mikhailov and O.Tarasov are grateful to the Physics Department of the Bielefeld University for warm hospitality. L.Avdeev and S.Mikhailov are thankful to the Volkswagen-Stiftung and O.Tarasov to the BMFT and RFFR grant for financial support. The authors are grateful to F.Jegerlehner for carefully reading the manuscript.

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