CROSS-SECTION ASYMMETRIES AROUND THE Z PEAK

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Abstract

A simple model-independent formula for cross-section asymmetries $A$ in fermion-pair production is derived, which may be used for the analysis of LEP 1 data,

$$A = \sum_n A_n (s - M_Z^2)^n.$$

The coefficient $A_0$ depends on the $Z$ boson exchange, $A_1$ additionally on the $\gamma Z$ interference, while the higher-order contributions are practically redundant. QED corrections are taken into account, and relations to other approaches are indicated.
1 Introduction

For fermion-pair production at LEP 1 energies,
\[ e^+e^- \to (\gamma, Z) \to f^+f^- (\gamma), \] (1)

experimentalists may present their cross-section asymmetries in a rather simple way:
\[ A(s) = A_0 + A_1 \left( \frac{s}{M_Z^2} - 1 \right) + A_2 \left( \frac{s}{M_Z^2} - 1 \right)^2 + \ldots \] (2)

Instead of \( s \), one can also use the centre-of-mass energy as a variable: \( \left( \frac{s}{M_Z^2} - 1 \right) = \left( \sqrt{s}/M_Z + 1 \right) \left( \sqrt{s}/M_Z - 1 \right) \sim 2(\sqrt{s}/M_Z - 1) \). Depending on the accuracy of the data, higher-order terms in the expansion may be neglected. Within \( \sqrt{s} = M_Z \pm \Gamma_Z \), the \( |(s/M_Z^2 - 1)| \) is less than 0.05, and coefficients beyond \( A_2 \) will hardly be within reach. With the data of the 1990 running period, for the \( \tau \)-polarization \( A_{pol} \), only the peak value \( A_0 \) has been determined, and the forward–backward asymmetry \( A_{FB} \) is known to be an almost linear function of \( s \). The LEP 1 data of 1992 will be much more precise and it seems to be reasonable to analyse the shapes of the cross-section asymmetries, and not only the peak values.

In this letter, it will be proved that (2) is a unique, model-independent ansatz for cross-section asymmetries around the \( Z \) peak. The coefficients \( A_n \) contain the complete physical information without approximations. Allowing for a smooth (and calculable) dependence of the \( A_n \) on \( s \), the complete QED corrections may be included in the ansatz (2). We derive (2) in two steps, first neglecting the photonic corrections.

2 Model-independent approach to asymmetries

A cross section may be parametrized in a unique way:
\[ \sigma(s) = \frac{4}{3} \pi \alpha^2 \left[ \frac{r_\gamma}{s} + \frac{sR + (s - M_Z^2)J}{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} + \sum_n \frac{r_n}{M_Z^2} \left( \frac{s}{M_Z^2} - 1 \right)^n \right]. \] (3)

For the total cross section \( \sigma_T \), a formula like (3) has been extracted from the Standard Model in [1]. In [2], this formula has been derived for \( \sigma_T \), but also for left–right \( (\sigma_{pol}, \sigma_{LR}) \) and forward–backward \( (\sigma_{FB}) \) cross-section differences from an ansatz for the scattering matrix. The definitions of mass and width are also discussed there. At LEP 1, the corresponding asymmetries are:
\[ A_{FB} = \frac{\sigma_{FB}}{\sigma_T}, \quad A_{pol} = \frac{\sigma_{pol}}{\sigma_T}. \] (4)

The four independent matrix elements \( \mathcal{M}_i \) for the scattering of helicity states,
\[ \mathcal{M}_i(s) = \frac{R_\gamma}{s} + \frac{R_{Z}^i}{s - s_Z} + \sum_n F_i^{(n)}(s_Z)(s - s_Z)^n, \] (5)
with \( s_Z = M_Z^2 - i M_Z \Gamma_Z \), are used to derive the cross sections. The \( R_{\gamma}, R_Z \), and \( F_i^{(n)} \) are complex constants, which characterize the scattering process. They compose the parameters \( r_{\gamma}, R, J, r_n \) in \([3]\) in \( \sigma_T, \sigma_{pol}, \sigma_{FB} \) \([2]\). For definiteness, we quote the quantum-mechanical interpretation of the coefficients:

\[
\begin{align*}
\kappa_T &= v_\gamma v_f, \\
\kappa_{FB} &= \frac{3}{2}\mu a_\gamma a_f, \\
\kappa_{pol} &= -v_\gamma a_\tau, \\
\kappa &= \frac{G_\mu M_Z^2}{\sqrt{2} 8\pi \alpha} = 0.3724 \left( \frac{M_Z}{91} \right)^2.
\end{align*}
\]

The parameter \( R \) in \([3]\) is built out of the residua of the \( Z \) pole in the four helicity amplitudes, the \( J \) comes from the \( \gamma Z \) interference, and \( \alpha(s) \) is the running QED coupling constant, which we assume to be known.

Explicit expressions for the coefficients \( R \) and \( J \) in \( \sigma_T, \sigma_{FB}, \sigma_{LR} \) and \( \sigma_{pol} \) in terms of the Standard Model have been derived in \([3]\). The \( r_n \) are additional, non-resonating quantum corrections and yield negligible contributions.

From \([3]\) and \([4]\), one easily derives expressions for the coefficients \( A_n \) in \([2]\). The peak value of the asymmetry is (index \( A = FB, pol \)):

\[
A_0 = \frac{R_A + \gamma^2 r_0^A}{R_T + \gamma^2 (r_0^T + r_\gamma^T)} \sim \frac{R_A}{R_T + \gamma^2 r_\gamma^T} \sim \frac{R_A}{R_T},
\]

where it is taken into account that \( r_\gamma^A = 0 \). Then, non-resonating quantum corrections \( \gamma^2 r_0 \sim \frac{2}{7} \gamma^2 \sim 2 \times 10^{-6} \) are neglected, and finally also the pure photonic contribution \( r_\gamma^T \), which is multiplied with \( \gamma^2 = \Gamma^2_Z/M_Z^2 \sim 0.75 \times 10^{-3} \). For the coefficient \( A_1 \), again after safely neglecting the quantum corrections \( r_0, r_1 \), one gets:

\[
A_1 = \left[ \frac{J_A}{R_A} - \frac{J_T}{R_T + \gamma^2 r_\gamma^T} + \frac{2\gamma^2 r_\gamma^T}{R_T + \gamma^2 r_\gamma^T} \right] A_0 \sim \left[ \frac{J_A}{R_A} - \frac{J_T}{R_T} \right] A_0.
\]
The higher-order coefficients are defined by a recurrence relation. Neglecting again the \( r_n \) and terms of order \( \mathcal{O}(\gamma^2) \),

\[
A_n = -\left(2 + \frac{J_T}{R_T}\right) A_{n-1} + \left[(1 + \frac{J_A}{R_A}) \delta_{n,2} - \left(1 + \frac{J_T + r_T}{R_T}\right)\right] A_{n-2}.
\] (15)

Assuming for a moment that photonic corrections may be neglected, and that the asymmetry (2) may be interpreted in terms of vector- and axial-vector couplings, it is:

\[
A_{FB}^0 = 3 \frac{a_e v_e a_f v_f}{(a_e^2 + v_e^2)(a_f^2 + v_f^2)}, \quad A_{pol}^0 = -\frac{2a_\tau v_\tau}{a_\tau^2 + v_\tau^2}.
\] (16)

A notation is used with \( a_f = 1, v_f = 1 - 4|Q_f|\sin^2\theta_W \). Further,

\[
A_{FB}^1 \sim \frac{3}{2\kappa} |Q_e Q_f| a_e a_f \frac{(a_e^2 + v_e^2)(a_f^2 + v_f^2) - 4v_e^2 v_f^2}{(a_e^2 + v_e^2)^2(a_f^2 + v_f^2)^2},
\] (17)

\[
A_{pol}^1 \sim -\frac{2}{\kappa} |Q_e Q_\tau| \frac{v_\tau a_\tau (a_e^2 - v_e^2)}{(a_e^2 + v_e^2)(a_\tau^2 + v_\tau^2)}.\] (18)

The \( A_0 \) is completely determined by the residua of the \( Z \) resonance. It is not influenced by the \( \gamma Z \) interferences, while the \( A_1 \) gets its leading contributions just from them. With a rising accuracy of the data, the photonic contributions may not be neglected in the above definitions \[4\]. In \( A_1 \), there is an additional suppression due to the small factor \((s/M_Z^2 - 1)\). With ideal data, the interpretation of the asymmetry is straightforward. The \( M_Z, \Gamma_Z, R_T \) and \( J_T \) may be determined from an analysis of the \( Z \) line shape \[1\]. Then, \( A_0 \) allows the determination of \( R_A \), and afterwards the \( \gamma Z \) interference \( J_A \) may be derived from \( A_1 \). Instead of \( R_A \), we may determine the ratio \( R_A/R_T \), which allows us to determine other coupling combinations than may be obtained from a line-shape analysis, thus improving determinations of the effective weak mixing angle. The problems connected with a secondary interpretation of the model-independent findings will not be discussed here, although they are important and interesting.

Since the \( A_0 \) is independent of the \( \gamma Z \) interferences, it is stable against several phenomena. For the \( Z \) peak position \( s_p \), one may derive the relation \[3\]:

\[
\Delta \sqrt{s_p} = \Delta M_Z + \frac{1}{4} \gamma^2 M_Z \Delta \left(\frac{J_T}{R_T}\right) + \ldots
\] (19)

between an uncertainty in \( M_Z \) and an uncertainty in the \( \gamma Z \) interference. The latter influences \( A_1 \). Similarly, for a hypothetical heavy gauge boson \( Z' \), the

\[1\] A serious line-shape analysis has at least four free parameters. While \( M_Z \) and \( \Gamma_Z \) are universal, the \( R_T, J_T \) are different for the different channels.
effects from its virtual exchange transform after a partial fraction decomposition into simple shifts of the $\gamma Z$ interferences [5]:

$$
\Delta \left( \frac{J_T}{R_T} \right) = -2 \frac{g^2}{g^2 M_{Z'}^2 - M_Z^2} \frac{(a_e a_e' + v_e v_e')(a_f a_f' + v_f v_f')}{(a_e^2 + v_e^2)(a_f^2 + v_f^2)}, \tag{20}
$$

and analogously for the cross-section differences. Again, the $A_1$ will be influenced, while $A_0$ is sensitive exclusively to the $ZZ'$ mixing effect. The $r_T^{\gamma}$ in the definition of $A_2$ is of the order of one, thus not suppressed by neglecting $\gamma$. As mentioned above, the whole $A_2$ is suppressed at LEP 1 with a factor of $(s/M_Z^2 - 1)^2 \leq 0.3\%$ and it contains no new physical information with respect to $A_{0,1}$.

### 3 Inclusion of photonic corrections

In this section, the modifications from the photonic corrections are discussed. Neglecting the initial–final interference bremsstrahlung, the cross sections $\sigma$ [eq. (3)] may be replaced by

$$
\bar{\sigma}(s) = \int dk \sigma(s') \rho(k), \tag{21}
$$

where $s' = (1-k)s$, and the radiator function $\rho(k)$ contains the QED corrections. An introduction to photonic corrections in the language of scattering amplitudes is given in [6]. A rather complete discussion of (21) may be found in [7, 8], and in the references therein. For the present purpose, after inserting (3) into (21), the cross section may be rewritten as follows:

$$
\bar{\sigma}(s) = \frac{4}{3} \pi \alpha^2 \left[ \frac{\bar{r}_\gamma}{s} + \frac{s \bar{R} + (s - M_Z^2) \bar{J}}{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} + \sum_n \frac{\bar{r}_n}{M_Z^2} \left( \frac{s}{M_Z^2} - 1 \right)^n \right]. \tag{22}
$$

Any of the barred parameters differs from its unbarred partner by a correction factor $C(s)$:

$$
\bar{R} = C_R(s) R, \quad C_R(s) = \mathcal{I} \left[ \frac{s' - M_Z^2}{s} \frac{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2}{(s' - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} \right], \tag{23}
$$

$$
\mathcal{I}[B] = \int dk B(s') \rho(k). \tag{24}
$$

The other correction factors are analogously defined:

$$
C_J(s) = \mathcal{I} \left[ \frac{s' - M_Z^2}{s - M_Z^2} \frac{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2}{(s' - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} \right], \tag{25}
$$

$$
C_r(s) = \mathcal{I} \left[ \frac{s}{s'} \right], \tag{26}
$$

$$
C_n(s) = \mathcal{I} \left[ \frac{(s' - M_Z^2)^n}{(s - M_Z^2)^n} \right]. \tag{27}
$$

These very small corrections may be properly taken into account [2]. Then, the number of free parameters increases.
The reader may wonder that some of the corrections seem to be singular at \( \sqrt{s} = M_Z \). This is not the case for the products \( C_{J,n}(s - M_Z^2) \). As may be seen from the corresponding definitions, these remain small compared with e.g. \( A_0 \), when \( \sqrt{s} \) approaches \( M_Z \). At this energy, the asymmetry is defined as the (smooth) limit from the neighbouring energies. In a more elegant notation, but unnecessarily sophisticated for applications, one could rewrite the asymmetry as a series in powers of \( (s - s_Z) \) as \( \bar{A}_n \), thus regularizing the \( C_{J,n} \). We remark that after inclusion of photonic corrections the parameters \( J \) and \( r_n \) may yield non-vanishing contributions at \( s = M_Z^2 \). This was not the case without the QED corrections, which smear out the effective energy.

The QED corrections are well-defined as soon as \( M_Z \) and \( \Gamma_Z \) are known. They are independent of the underlying dynamics of the scattering process. It is not difficult to collect explicit expressions for the \( C \) functions from the literature. They, of course, depend on the handling of the photonic phase space, the inclusion of higher orders, and on acceptance cuts. For example, the initial-state corrections in \( A_0 \) (with possible inclusion of soft photon exponentiation), with a cut on the energy of the emitted photon, are \( C_T^R = R_T(1,1), C_F^B = R_{FB}(1,1) \), where the \( R_{T,FB} \) are defined in eqs. (56)–(57) and (78) of \( \text{[8]} \). In the simplest case [initial state radiation to \( O(\alpha) \) without cuts], it is \( C_T^R = 1 + \frac{\alpha}{\pi} H_0^T, C_F^B = 1 + \frac{4}{3} \frac{\alpha}{\pi} H^T_3 \), with the \( H_{0,3} \) to be taken from \( \text{[9]} \).

Taking into account the QED corrections, the experimental data may be fitted with the ansatz

\[
\bar{A}(s) = \bar{A}_0 + \bar{A}_1 \left( \frac{s}{M_Z^2} - 1 \right) + \bar{A}_2 \left( \frac{s}{M_Z^2} - 1 \right)^2 + \ldots \tag{28}
\]

The \( \bar{A}_n \) may be obtained from the \( A_n \), replacing everywhere in the corresponding definitions the unbarrred variables by barred ones. For the leading contribution to the forward–backward asymmetry, the explicit expression is:

\[
\bar{A}_{F B}^0 = \frac{C_{F B}^R(s)}{C^T_R(s)} \frac{R_{F B}}{R_T + [C^T_T(s)/C^T_R(s)] \gamma^2 r_T^T} \sim 0.998 \frac{R_{F B}}{R_T + 0.001}. \tag{29}
\]

Note that the radiator function \( \rho_{FB}(k) \) in (24), which must be used for the calculation of \( \sigma_{FB} \), differs from the radiator \( \rho_T(k) \). The latter is used both for \( \sigma_T \) and \( \sigma_{pol} \), and the expression for \( A_{pol} \) simplifies correspondingly. The leading term is:

\[
\bar{A}_{pol}^0 = \frac{R_{pol}}{R_T + [C^T_T(s)/C^T_R(s)] \gamma^2 r_T^T} \sim \frac{R_{pol}}{R_T + 0.001}. \tag{30}
\]

Further, neglecting the strongly suppressed contributions to \( A_1 \) (index \( A = FB, pol \)):

\[
\bar{A}_1 = \frac{C^T_T(s)}{C^T_R(s)} \left[ \frac{J_A}{R_A} - \frac{J_T}{R_T} \right] \bar{A}_0. \tag{31}
\]
The explicit numerical values in (29) and (30) may be taken from Table 1, which is calculated with the FORTRAN package ZFITTER \[8, 10, 11\].

\[
\sqrt{s} \quad M_Z - 2\Gamma_Z \quad M_Z - \Gamma_Z \quad M_Z - \frac{1}{2}\Gamma_Z \quad M_Z \quad M_Z + \frac{1}{2}\Gamma_Z \quad M_Z + \Gamma_Z \quad M_Z + 2\Gamma_Z \\
C_{FB}^T \quad 0.7784 \quad 0.7331 \quad 0.7078 \quad 0.7350 \quad 0.9367 \quad 1.2209 \quad 1.8120 \\
C_{FB}^T / C_{R}^T \quad 0.9977 \quad 0.9980 \quad 0.9982 \quad 0.9982 \quad 0.9981 \quad 0.9978 \quad 0.9964 \\
C_{T}^T / C_{R}^T \quad 1.7422 \quad 1.8565 \quad 1.9264 \quad 1.8582 \quad 1.4601 \quad 1.1215 \quad 1.7569 \\
C_{FB}^J / C_{FB}^T \quad 1.0881 \quad 1.1111 \quad 1.1592 \quad [-0.004] \quad 0.6176 \quad 0.5649 \quad 0.4384 \\
C_{T}^J / C_{R}^T \quad 1.0905 \quad 1.1130 \quad 1.1610 \quad [-0.004] \quad 0.6159 \quad 0.5631 \quad 0.4358 \\
\]

Table 1: QED corrections to the parameters of the model-independent asymmetry formulae; \( M_Z = 91.146 \) GeV, \( \Gamma_Z = 2.499 \) GeV.

For this purpose, one should use the branch which relies on the S-matrix ansatz \[12\]. The maximal acollinearity of the final-state fermions is assumed to be \( \xi = 15^\circ \), and the minimal energy of one of the fermions to be \( E_{\text{min}} = 20 \) GeV (standard cuts of ZFITTER with flag ZUCUTS=1). The photonic corrections to the asymmetries are remarkably stable against a variation of the cut conditions. Higher-order corrections with common exponentiation of initial- and final-state corrections are taken into account.

The photonic corrections to \( A_0 \) are nearly negligible. The reason is that the corrections \( C_T \) and \( C_{FB} \) differ only owing to hard-photon emission, which is strongly suppressed at the \( Z \) peak \[13\]. Explicit expressions for their ratio may be found in eqs. (81)–(86) of \[8\]. The photonic corrections to \( A_1 \) show a completely different behaviour. This is due to the ratio \( C_J(s)/C_R(s) \). In (31) we took into account that this ratio is nearly identical for \( \sigma_{FB} \) and \( \sigma_{T,pol} \). As has been discussed in \[8, 14\], there is an essential difference between the two corrections \( C_J(s) \) and \( C_R(s) \): while the pure \( Z \) exchange cross section (i.e. \( C_R \)) develops a radiative tail, the \( \gamma Z \) interference (i.e. \( C_J \)) does not. Consequently, their ratio is smooth and of order one below the resonance, while above it becomes considerably smaller since \( C_R(s) \) grows up. Mainly for this reason, the measured asymmetries are nearly linear functions of \( \sqrt{s} \) at \( \sqrt{s} < M_Z \), and become suppressed beyond the peak. In principle, the radiative tail may be avoided by a cut on the allowed energy of the emitted photons \[14\]:

\[
\frac{E_\gamma}{E_{\text{beam}}} < \Delta = 1 - \frac{M_Z^2}{s}.
\]
At LEP 1, where $s$ is near to $M_Z^2$, this condition is rather restrictive; e.g. at $\sqrt{s} = M_Z + 2\Gamma_Z$, it is $\Delta = 0.1$. Thus, usually one presents data including radiative corrections (see e.g. figures 19a-c of [15]). In the present approach, the ratio $C_J/C_R$ is the only QED correction, which is essentially energy dependent. As has been mentioned above, near $\sqrt{s} = M_Z$, one should better enumerate the smooth product $(s/M_Z^2 - 1)C_J/C_R$. This has been done in Table 1; see the numbers in square brackets there.

The higher-order coefficients $A_n$ are composed out of the two first ones, and the same is true for their photonic corrections.

4 Discussion

From the model-independent $Z$ line-shape formulae, we derived the corresponding expressions for the forward–backward asymmetry $A_{FB}$ and the $\tau$ polarization $A_{pol}$ at LEP 1 energies. The analytic expressions are valid for the leptonic and $b$-quark forward–backward asymmetries as well as for the $\tau$ polarization. The remarkably simple power series in $(s - M_Z^2)$ may cover in their coefficients the photonic corrections as complete as the line-shape formulae do. The asymmetries are defined by only two free parameters $A_0$ and $A_1/A_0$ or, alternatively, $R_A, J_A$ ($A = FB, pol$). The latter ones are related to the $ZZ$ and $\gamma Z$ contributions to the corresponding asymmetric cross-section combinations. In a subsequent step, one may determine effective couplings or radiatively corrected Standard Model parameters from the measured model-independent numbers. Additional, non-resonating quantum corrections may be neglected. The photonic corrections are defined such that they depend exclusively on $s, M_Z, \Gamma_Z$. The FORTRAN package ZFIT ER may be used for their calculation. For $A_0$ they are extremely small. The corrections to $A_1$ are dominated by the radiative tail.

From a combined analysis of the line shape and of asymmetries, one may try to determine the basic quantities of the S-matrix approach, i.e. the four complex residua of the $Z$-boson pole $R^i_Z$ in (5). For a given channel, this deserves the measurement of four independent sets of parameters $R, J$. For lepton-pair production, the three energy-dependent quantities $\sigma_T, A_{FB}, A_{pol}$ have been determined experimentally. These are sufficient for at least the determination of the real parts of the leptonic $R^i_Z$. As long as there is no beam polarization available at LEP 1, one could try to measure as a fourth, independent leptonic observable the $\tau$ polarization in forward direction, $\lambda^F_T$, or the asymmetry $\lambda^{FB}_T$, as was proposed in [11, 16].
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