Second order contributions to elastic large-angle Bhabha scattering *

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Abstract: We derive the coefficient of the $\mathcal{O}(\alpha^2 \log(s/m_e^2))$ fixed order contribution to elastic large-angle Bhabha scattering. We adapt the classification of infrared divergences, that was recently developed within dimensional regularization, and apply it to the regularization scheme with a massive photon and electron.

Keywords: QED, Bhabha scattering, NNLO Computations, infrared and collinear divergences.

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1. Introduction

Bhabha scattering has been measured with great accuracy by the experiments at the LEP and SLC colliders. This has led to many works by numerous authors being devoted to calculating the theoretical cross section for this process as accurately as possible [1, 2]. Yet, the current accuracy of theoretical predictions for large angle Bhabha scattering at LEP2 is still not completely satisfactory [3], and limits the bounds that can be set on some kinds of potential new physics effects, for instance, those coming from large extra dimensions [4].

Future linear electron positron colliders, e.g. the proposed TESLA accelerator, will not have monochromatic beams because the electrons and positrons can emit beamstrahlung before they collide. Measurement of the resulting acollinearity angle in large angle Bhabha events provides a way to determine the luminosity spectrum [5, 6].

Large angle Bhabha scattering is also important at electron-positron colliders running at centre of mass energies of a few GeV, such as BEPC, VEPP-2M, DAPHNE, and the B-factories PEP-II and KEK-B, where it is used to measure the integrated luminosity [7]. At present, the theoretical uncertainty on the differential cross section of this process is one of the limiting factors on the precision of the luminosity determination [8].

In order to improve the theoretical predictions, it is necessary to include higher order radiative corrections. In this paper, we consider only pure QED corrections. These fall in different orders of magnitude $\alpha L$, $\alpha^2 L^2$, $\alpha^2 L$, $\alpha^3 L^3$, where $\alpha$ is the fine-structure constant and $L = \log(s/m_e^2)$. The large logarithm $L$ is related to collinear divergences that would appear if the electron mass $m_e$ were zero. The $O(\alpha^2 L)$ terms are so far only partially known. Arbuzov, Kuraev and Shaikhatdenov have calculated the contributions from soft one- and two-photon bremsstrahlung, squared one-loop graphs, and the interference between two-loop vertex graphs and tree level terms [9]. However, they did not calculate the contribution due to two-loop box graphs, which has to be included to complete the $O(\alpha^2 L)$ correction.

In principle, it should be possible to extract the missing two-loop box contribution from the work of Bern, Dixon and Ghinculov [10], who have presented a complete formula for the two-loop virtual corrections to Bhabha scattering. However, these authors set the electron mass to zero and regularized all infrared and collinear divergences dimensionally, with $d = 4 - 2\epsilon$, unlike the authors of ref [9], who used the traditional method of an electron and a photon mass. For this reason, the results obtained by the two groups cannot be directly combined.

If one is only interested in the logarithmically enhanced term, one does not actually need the full result of ref. [10], but only the terms containing powers of $1/\epsilon$. The correspondence between poles in $\epsilon$ of the dimensional regularization scheme and logarithms $\log(m_e)$ and $\log(m_\gamma)$ of the massive regulator scheme has been worked
out in detail at the one-loop level [11]. However, it is not a priori clear how to extend this to two loops. Therefore, in this paper, we take a slightly different approach. We rely on an elegant classification, proposed by Catani, of the singularities of on-shell two-loop amplitudes in QCD [12]. His results were confirmed for the QED case by [10]. Catani’s formalism was developed using dimensional regularization. However, it ought to be valid in both regularization schemes. We shall translate this formalism into the massive regulator scheme, and then use it to predict the $O(\alpha^2 L)$ contribution to the cross section for large angle Bhabha scattering. We limit ourselves to graphs without photon vacuum polarization insertions, since those can be treated as a separate class [13].

Our result should be directly applicable to the energy range of a few GeV, where non-QED effects, from graphs involving $Z$-boson exchange, are still very small. There are several Monte Carlo programs that were specifically designed for this energy range. These include the programs BABAYAGA and LABSPV [7] and LABSMC [14]. (Monte Carlo programs written for LEP1 and LEP2 are reviewed in refs. [2, 3]). The precision of these programs depends on the energy and on the details of the cuts applied. In the case of LABSMC, under typical conditions, the uncertainty is estimated by the authors to be around 0.2% [14], mainly due to missing terms of order $O(\alpha^2 L)$. The result obtained in this paper will allow this uncertainty to be reduced.

2. Infrared factorisation formulae

A generic (renormalized) QED matrix element can be expanded as a series in $\alpha$ as follows,

$$|\mathcal{M}\rangle = \left(\frac{\alpha}{2\pi}\right)^n \left( |\mathcal{M}^{(0)}\rangle + \left(\frac{\alpha}{2\pi}\right) |\mathcal{M}^{(1)}\rangle + \left(\frac{\alpha}{2\pi}\right)^2 |\mathcal{M}^{(2)}\rangle + O(\alpha^3) \right)$$

where $|\mathcal{M}^{(i)}\rangle$ represents the $i$-loop contribution and where the overall power $n$ may be half-integer. According to Catani [12], these amplitudes obey the following factorization formulae,

$$|\mathcal{M}^{(1)}\rangle = I^{(1)} |\mathcal{M}^{(0)}\rangle + |\mathcal{M}^{(1),fin}\rangle$$
$$|\mathcal{M}^{(2)}\rangle = I^{(2)} |\mathcal{M}^{(0)}\rangle + I^{(1)} |\mathcal{M}^{(1)}\rangle + |\mathcal{M}^{(2),fin}\rangle$$

where $I^{(1)}$ and $I^{(2)}$ are operators that depend on the regularization scheme and contain all of the singularities of the infrared regulator. The remainders $|\mathcal{M}^{(1),fin}\rangle$ and $|\mathcal{M}^{(2),fin}\rangle$ are finite. In dimensional regularization, for QED processes with massless electrons,

$$I^{(1)}(\epsilon) = \frac{1}{2\Gamma(1-\epsilon)} \sum_{i=1}^{n} \sum_{i\neq j} \epsilon_i \epsilon_j \left( \frac{1}{\epsilon^2} + \frac{3}{2\epsilon} \right) \left( \frac{\mu^2 e^{-i\lambda_{ij}\pi}}{2p_i \cdot p_j} \right)^\epsilon$$

(2.3)
where $\lambda_{ij} = -1$ if $i$ and $j$ are both incoming or outgoing partons and $\lambda_{ij} = 0$ otherwise, and $e_i$ is the electric charge (minus the electric charge) of an outgoing (incoming) radiating particle with momentum $p_i$. Eq. (2.3) is obtained from the QCD result [12] with the replacements $C_A \rightarrow 0$, $C_F \rightarrow 1$, $T_R \rightarrow 1$ and $T_i \cdot T_j \rightarrow e_i e_j$.

Because we exclude graphs with photon vacuum polarization insertions, we neglect terms proportional to the $\beta$ function. The operator $I^{(2)}$ is then given by

$$I^{(2)} = -\frac{1}{2} I^{(1)} I^{(1)} + H^{(2)},$$

where $H^{(2)}$ contains only single poles and has the form

$$H^{(2)} = \frac{1}{4\epsilon} \frac{e^{i\gamma}}{\Gamma(1-\epsilon)} \sum_{i=1}^{n} \sum_{i \neq j}^{n} e_i e_j \left( \frac{\mu^2 e^{-i\lambda_{ij}\pi}}{2p_i \cdot p_j} \right)^{2\epsilon} H^{(2)},$$

where

$$H^{(2)} = \frac{3}{8} - 3\zeta_2 + 6\zeta_3. \tag{2.6}$$

Ref. [12] does not give a general formula for $H^{(2)}$, and at present it is necessary to derive it from an explicit two-loop calculation, such as the two-loop QCD computation of the electromagnetic form factor of the quark [15]. We note that in Ref. [12] the summation over radiating pairs $i$ and $j$ is not explicitly included in the definition of $H^{(2)}$. However, it has been found to reproduce correctly the results of explicit calculations in QED [10] and QCD [16]. We also note that the terms of $O(1)$ in $H^{(2)}$ are presently a matter of choice and can be altered by a redefinition of $|M^{(2),\text{fin}}\rangle$.

3. The massive photon and massive electron regularisation scheme

The factorization formulae (2.2) should hold for any choice of infrared regulator. Therefore, we can find a translation between schemes by comparing explicit calculations. For example, the one-loop correction to the electron photon vertex (see, e.g. Ref. [17]), serves to fix $I^{(1)}$ in the scheme where the photon has mass $\lambda$ and the electron has mass $m$. We write the Dirac form factor $F_1(q^2)$ of the electron photon vertex

$$\Gamma_{\mu} = \gamma_{\mu} F_1(q^2) + \frac{i}{2m} \sigma_{\mu
u} q^\nu F_2(q^2)$$

as

$$F_1(q^2) = 1 + \frac{\alpha}{\pi} F_1^{(1)}(q^2) + \left( \frac{\alpha}{\pi} \right)^2 F_1^{(2)}(q^2) + O\left( \alpha^3 \right). \tag{3.2}$$

\footnote{In order to follow more closely the notation of Ref. [9], we expand in $\alpha/\pi$ rather than $\alpha/(2\pi)$ in the remainder of this paper. Therefore, the definitions of $I^{(1)}$ and $H^{(2)}$ differ by factors of 1/2 and 1/4 respectively from those in the previous section.}
By equating
\[ F_1^{(1)}(q^2) = I^{(1)}(q^2) + F_1^{(1),\text{fin}}(q^2), \] (3.3)
we see that
\[ I^{(1)}(q^2) = (L - 1) \left( \frac{1}{2} L_\lambda + 1 \right) - \frac{1}{4} L^2 - \frac{1}{4} L + \frac{\pi^2}{12}, \] (3.4)
where \( L_\lambda = \log \left( \frac{\lambda^2}{m^2} \right) \), \( L = \log \left( -\frac{q^2}{m^2} \right) \) if \( q^2 < 0 \) and \( L = \log \left( \frac{q^2}{m^2} \right) - i\pi \) if \( q^2 > 0 \). There is a possible ambiguity in assigning the constant pieces to \( I^{(1)} \) or \( F_1^{(1),\text{fin}} \). Our choice corresponds to \( F_1^{(1),\text{fin}} = 0 \). Similarly, by examining the two-loop vertex, we can fix \( H^{(2)} \). To do this, we take the large scale limit of the \( \mathcal{O}(e^4) \) two-loop vertex correction \( F_1^{(2)}(q^2) \) computed by Barbieri, Mignaco and Remiddi [17] (without the contribution from the vacuum polarization graph). Up to terms that do not depend on the regulators, we expect that,
\[ F_1^{(2)}(q^2) = \frac{1}{2} I^{(1)}(q^2) I^{(1)}(q^2) + H^{(2)}(q^2) + \mathcal{O}(1). \] (3.5)

We find that \( H^{(2)}(q^2) \) is proportional to a single large logarithm
\[ H^{(2)}(q^2) = \frac{1}{4} L H^{(2)} \] (3.6)
where \( H^{(2)} \) is given by Eq. (2.6). In fact, one might wonder whether changing the constants in \( I^{(1)}(q^2) \) would give rise to a different \( H^{(2)} \). This is not the case since any change is absorbed by the necessary alteration to \( F_1^{(1),\text{fin}} \).

Armed with these operators using mass regularization, we can compute the large logarithmic corrections to the two-loop contribution to Bhabha scattering keeping the electron mass and using a small photon mass as the infrared regulator. For \( 2 \to 2 \) scattering, there are six radiating pairs, so that
\[ I^{(1)} = 2 \left( I^{(1)}(s) + I^{(1)}(t) - I^{(1)}(u) \right), \] (3.7)
and
\[ H^{(2)} = 2 \left( H^{(2)}(s) + H^{(2)}(t) - H^{(2)}(u) \right). \] (3.8)

The second order virtual contribution to Bhabha scattering comes from the square of the one-loop graphs and the interference of tree and two-loop graphs. We can write this contribution in factorized form as
\[ \frac{d\sigma^{VV}}{d\sigma_0} = \left( \frac{\alpha}{\pi} \right)^2 \Delta_{VV}, \] (3.9)
where the lowest order differential cross section is given by
\[ d\sigma_0 = \frac{\alpha^2}{s} \left( \frac{1 - x + x^2}{x} \right)^2 d\Omega, \] (3.10)
with \( x = -t/s \) and \( 1 - x = -u/s \). Up to constant terms,

\[
\Delta_{VV} = -\frac{1}{2} \left( I^{(1)} + I^{(1)*} \right)^2 + \delta_V \left( I^{(1)} + I^{(1)*} \right) + \left( H^{(2)} + H^{(2)*} \right)
\]

(3.11)

with \( I^{(1)} \) and \( H^{(2)} \) given by Eqs. (3.7) and (3.8) respectively. The second term involves the one-loop virtual contribution \( \delta_V \) (see, e.g. [18, 19] and references therein).

In the notation of ref. [9], it is given by,

\[
\delta_V = 4 \log \frac{m}{\lambda} \left( 1 - L + \log \left( \frac{1-x}{x} \right) \right) - L^2 + 2L \log \left( \frac{1-x}{x} \right) - \log^2(x) \\
+ \log^2(1-x) + 3L - 4 + f(x),
\]

(3.12)

with\(^2\)

\[
f(x) = (1 - x + x^2)^{-2} \left[ \frac{\pi^2}{12} (4 - 8x + 27x^2 - 26x^3 + 16x^4) \\
+ \frac{1}{2} (-2 + 5x - 7x^2 + 5x^3 - 2x^4) \log^2(1-x) + \frac{1}{4} x (3 - x - 3x^2 + 4x^3) \log^2(x) \\
+ \frac{1}{2} (6 - 8x + 9x^2 - 3x^3) \log(x) - \frac{1}{2} x (1 + x^2) \log(1-x) \\
+ \frac{1}{2} (4 - 8x + 7x^2 - 2x^3) \log(x) \log(1-x) \right].
\]

(3.13)

Expanding \( \Delta_{VV} \) yields all terms containing at least one power of the large logarithm \( L = \log(s/m^2) \) as well as all logarithms of the photon mass regulator. We have checked that in the small angle limit, \( x \to 0 \), to the logarithmic accuracy we are working at, Eq. (3.11) reduces to the known result [20]

\[
\Delta_{VV} = 6 \left( F_1^{(1)}(t) \right)^2 + 4 F_1^{(2)}(t),
\]

(3.14)

which follows from a generalized eikonal representation [21] for the Bhabha scattering amplitude for small angles.

### 4. The scattering cross-section at \( \mathcal{O}(\alpha^2) \)

To make a physical prediction, the double virtual contribution must be combined with the one-loop contribution with single soft emission and the tree level double soft emission. The second order correction to the one-loop virtual photon emission corrected cross section, due to the emission of a single real soft photon having energy less than \( \Delta \varepsilon \), can be written down in the factorized form [9]

\[
\frac{d\sigma^{SV}}{d\sigma_0} = \frac{\alpha}{\pi} \delta_S \frac{\alpha}{\pi} \delta_V = \left( \frac{\alpha}{\pi} \right)^2 \Delta_{SV},
\]

(4.1)

\(^2\)There are some misprints in the formula for \( f(x) \) in ref. [9].
where $\delta_V$ is given in Eq. (3.12) and

$$
\delta S = 4 \log \left( \frac{m \Delta \varepsilon}{\lambda \varepsilon} \right) \left( L - 1 + \log \left( \frac{x}{1 - x} \right) \right) + L^2 + 2L \log \left( \frac{x}{1 - x} \right) + \log^2(x)
- \log^2(1 - x) - \frac{2\pi^2}{3} + 2 \text{Li}_2(1 - x) - 2 \text{Li}_2(x).
$$

(4.2)

Here, $\varepsilon = \sqrt{s/4}$ is the energy of the electron and positron beams. The dilogarithm function is defined as

$$
\text{Li}_2(x) = - \int_0^x \frac{dy}{y} \log(1 - y).
$$

(4.3)

The contribution from two independently emitted soft photons each with energy $\omega_1, \omega_2 \leq \Delta \varepsilon$ is given by [9]

$$
\frac{d\sigma_{SS}}{d\sigma_0} = \frac{1}{2} \left( \frac{\alpha}{\pi} \right)^2 \delta_S^2 \equiv \left( \frac{\alpha}{\pi} \right)^2 \Delta_{SS},
$$

(4.4)

where the statistical factor $1/2!$ is due to the identity of photons.

The combination $\Delta_{SS} + \Delta_{SV}$ is written in expanded form in Ref. [9] where terms like $L^4, L^3 L_\lambda, L^2, L^2 L_\lambda, LL_\lambda$ and $LL_\lambda^2$ are produced. All of these terms precisely cancel against similar terms in $\Delta_{VV}$. The final second order contribution to the cross section through to $O(1)$ is given by

$$
\frac{d\sigma}{d\sigma_0} = \left( \frac{\alpha}{\pi} \right)^2 (\Delta_{VV} + \Delta_{SV} + \Delta_{SS})
= \left( \frac{\alpha}{\pi} \right)^2 \left[ L^2 \left( 8 \log^2 \left( \frac{\Delta \varepsilon}{\varepsilon} \right) + 12 \log \left( \frac{\Delta \varepsilon}{\varepsilon} \right) + \frac{9}{2} \right)
+ L \left( A \log^2 \left( \frac{\Delta \varepsilon}{\varepsilon} \right) + B \log \left( \frac{\Delta \varepsilon}{\varepsilon} \right) + C \right) \right].
$$

(4.5)

The single logarithmic coefficients $A$, $B$ and $C$ are given by,

$$
A = 16 \log \left( \frac{x}{1 - x} \right) - 16,
$$

(4.6)

$$
B = 8 \text{Li}_2(1 - x) - 8 \text{Li}_2(x) + 12 \log \left( \frac{x}{1 - x} \right) + 4f(x) - 28 - \frac{8}{3} \pi^2,
$$

(4.7)

$$
C = 6 \text{Li}_2(1 - x) - 6 \text{Li}_2(x) + 3f(x) + 6\zeta_3 - \frac{93}{8} - \frac{5}{2} \pi^2.
$$

(4.8)

Coefficients $A$ and $B$ agree with those obtained by Arbuzov et al. [9], while $C$ is the main new result of the present work.

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3up to a slight misprint in the expression for $B$, denoted there by $z_1.$
5. Summary

In this form Eq. (4.5) cannot be used to compare directly with experiment. The reason is that an experimental set-up involves a complicated detector that is not represented by the simple energy cuts on the photons which we used here. Such effects have to be modeled by a Monte-Carlo calculation. However any uncertainty in the cross section is now reduced to $O(\alpha^2)$ terms without enhanced logarithms. Given the final simplicity of the method one can wonder whether it is possible to use the results of [10] to extract the finite parts of the cross section. At present this appears to be not simply possible, as we have no way to tell whether the $1/\epsilon$ poles contained in $H^{(2)}$ (Eq. (2.5)) correspond to $L$ or for instance $L - 1$. Presumably it is possible to establish this connection by further expanding the vertex functions. The alternative would be to calculate the box graphs within the mass regulator scheme. However this is quite a challenge. For application to analogous processes in QCD we remark however that this last uncertainty plays no role as it gets absorbed in the definition of the structure functions.

The calculation presented in this paper removes a major obstacle to improvement of the precision of Monte Carlo programs for large angle Bhabha scattering at centre of mass energies of a few GeV [14]. The result can also be used at higher energies, provided that additional contributions due to $Z$-exchange diagrams are included. Whether it is possible to obtain them by an extension of the method we have used here is a question that requires further investigation.

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