EPJ Web of Conferences **125**, 04005 (2016) DOI: 10.1051/epjconf/201612504005
QUARKS-2016

Missed radiative corrections in muon g-2 and proton charge radius measurements

Andrej Arbuzov[1,3,⋆] and *Tatiana* Kopylova[2]

[1] *Bogoliubov Laboratory for Theoretical Physics, JINR, Dubna 141980, Russia*
[2] *Dubna State University, Dubna 141982, Russia*

Abstract. QED radiative corrections to the muon anomalous magnetic moment and elastic electron-proton scattering are discussed. It is shown that a collective effect due to mutual interaction of muons within experimental conditions might provide a contribution to the observed muon magnetic moment. This effect is parameterized by an effective mean shift of muons off their mass shells. Higher order corrections to elastic electron-proton at low energies are systematically treated within the leading and next-to-leading logarithmic approximation. The corrections are relevant for the modern experiments on proton form factor and charge radius definition.

1 Introduction

This report is devoted to discussion of effects due to higher order QED radiative corrections (RC) in observables which are being studied in modern low-energy experiments. A very high precision is achieved some experiments, that makes a challenge for construction of adequately accurate theoretical predictions. In particular various effects of radiative corrections have to be re-considered. By *missed* RC we mean effects which were either not taken into account or treated in a wrong way. The report is based on two papers [1] and [2].

2 Off-mass-shell muon anomalous magnetic moment

We have [3] the following difference between the experimental data and theoretical predictions for the anomalous magnetic moment of muon:

$$\Delta a_\mu^{\mathrm{exp-SM}} = 288(63)(49)\cdot 10^{-11} \qquad (1)$$

It is almost twice as large as the pure weak contribution $\Delta a_\mu^{\mathrm{weak}} = 153.6(1.0)\cdot 10^{-11}$. An explanation of the deviation by a contribution due to some effects beyond the SM, requires introduction of a new energy scale being rather close to the EW one, while ongoing searches (in particular at LHC) more and more disfavor finding new physics at such a scale. Attempts to explain the difference by some effect of strong interactions result in continuous efforts in calculations of the corresponding non-perturbative contributions and trying to fix them using relations to experimentally observables quantities, see *e.g.*

⋆ e-mail: arbuzov@theor.jinr.ru

review [4]. Presently, there remains some valuable uncertainty in QCD effects, but the possibility to explain in this way the difference (1) is very unlikely. Another possibility could be an error in the experimental analysis. It will be verified by the forthcoming experiments Muon $g-2$ [5] at Fermilab (USA) and J-PARC $g-2$ [6] at KEK (Japan).

In order to resolve the puzzle, we have to look for the whole spectrum of possibilities. Here we discuss the scenario that some additional (but standard) interactions of muons within given experimental conditions contribute to the observed gyromagnetic ratio. Obviously interactions of a muon with his neighbors in the beam can lead to a (small) effective shift of its mass shell. Such collective effects of electromagnetic interactions between charged particles become large for intensive beams [7]. In the $g-2$ experiment the beam was not dense but the effect we are looking for is small. Remind also a weakened conservation of energy and momentum required to describe oscillations of atmospheric neutrinos originated from muon (and pion) decays.

The calculations of the one-loop fermion from factor with one external fermion line being off-shell was presented in the book [8]. Here we will consider only the off-diagonal situation (one fermion off-shell and the other one on-shell).

We choose the following kinematics:

$$k^2 = 0, \qquad p^2 = -m^2, \qquad (p+k)^2 + m^2 = 2pk = \kappa m^2, \tag{2}$$

where k is the photon 4-momentum; p and $p+k$ are the electron momenta. One of them is off shell. Note that metric $(-,+,+,+)$ is used. Dimensionless parameter κ describes the degree of the off-shellness, $\kappa < 0$. We will consider small values $|\kappa| \ll 1$.

At the one-loop level, the fermion-photon vertex receives contributions from scalar, vector and tensor Feynman integrals J_0, J_σ, $J_{\sigma\tau}$. The scalar and the tensor integrals are infrared and ultraviolet divergent, respectively. But as it is known from the standard calculations of $(g_f - 2)$, just the vector integral contributes to this quantity. For the off-mass-shell case it reads [8]

$$\frac{m^2}{i\pi^2} J_\sigma = \Big(J_0 - \frac{\ln|\kappa|}{\kappa - 1}\Big) p_\sigma + \Big(2J_0 - \frac{\kappa-2}{\kappa-1}\ln|\kappa| - 2\Big)\frac{k_\sigma}{\kappa},$$
$$J_0 = \frac{1}{\kappa}\Big[\mathrm{Li}_2(1) - \mathrm{Li}_2(1-\kappa)\Big], \qquad \mathrm{Li}_2(x) \equiv -\int\limits_0^x \frac{\ln(1-y)}{y} dy. \tag{3}$$

In the limit $\kappa \to 0$ and $k \to 0$, the vector integral leads to the well know one-loop result $\Delta a_f^{(1)} = \alpha/(2\pi)$ received first by J. Schwinger [9]. Expanding in κ gives the first correction due to the off-shellness:

$$\Delta a_f^{(1,\kappa)} = \frac{\alpha}{2\pi}\Big[1 + \delta a_f^{(\kappa)}\Big], \qquad \delta a_f^{(\kappa)} = \Big(\frac{1}{4} + \frac{\ln|\kappa|}{2}\Big)\kappa + O(\kappa^2). \tag{4}$$

Assuming that the difference (1) is due to the off-shellness effect, we get the equation to define the value of κ:

$$\Delta a_\mu^{\mathrm{exp-SM}} \approx 3 \cdot 10^{-9} = \frac{\alpha}{2\pi}\delta a_\mu^{(\kappa)}. \tag{5}$$

Numerical solution gives $\kappa \approx -3.5 \cdot 10^{-7}$. Note that for small values of κ the shift has the proper sign and the solution exists. Such a value corresponds to off-shellness of a muon of the order $m|\kappa| \sim 35$ eV. One can see that such an off-shellness is very small compared to the muon mass. But to get it due to an external electric field one requires a rather dense muon beam which is not the case of the former Brookhaven $g-2$ experiment. The new experiments at Fermilab and J-PARC will have quite different experimental condition. This will allow to test indirectly the potential collective effect. It is well known that collective effects are important in experiments with high-intensity beams. We claim such effects can also contribute to observables which are measured at an extremely high precision.

3 Higher order radiative corrections to elastic electron-proton scattering

Here we will discuss the treatment of QED radiative corrections in the Mainz Microtron (MAMI) experiment [10]. Elastic electron-proton scattering was measured there with point-to-point errors of the order of a few permille. That allowed to extract information on electric and magnetic form factors of the proton with high precision. Extrapolation of the electric form factor to the zero momentum transfer provided the value of the proton charge radius. The latter appeared being inconsistent with the value extracted from the muonic hydrogen spectrum [11]. This fact stimulates verification of all relevant elements of the studies. In particular, the effects of radiative corrections have to be treated with care.

We consider the process

$$e(p_1) + p(P_1) \longrightarrow e(p_2) + p(P_2) + (n\gamma,\ e^+e^-). \tag{6}$$

The initial electron energy $E_1 = p_1^0 \equiv E$ is of the order 1 GeV, $E \gg m_e$. The momentum transfer squared $Q^2 = -(p_2 - p_1)^2$ lies the range $0.003 < Q^2 < 1$ GeV^2. The condition $Q^2 \gg m_e^2$ holds for the whole range.

One-loop QED corrections to this process are well known. They are naturally separated into the following parts: i) real and virtual corrections to the electron line, ii) real and virtual corrections to the proton line, iii) interference of amplitudes of the first two types, iv) the effect due to vacuum polarization. The corresponding analytic results were reproduced in [10]. Among one-loop corrections, there is still an open discussion about the proper treatment of double photon exchange contributions, see *e.g.* papers [12, 13] and references therein. We agree with the importance of this point, but it goes beyond the scope of our present study.

To estimate the numerical effect of radiative corrections one has to take into account concrete experimental conditions. Of course, to get the final answer one should include the corrections into the whole program of the data analysis. But our task here will be just to present analytic results with simple estimates of their impact. So we will simplify the set-up (still following the main features of the experiment):
— we assume that the measurement is based on the detection of the final electron energy and momentum,
— the electron is detected “bare”, *i.e.* without possible accompanying photons,
— there is just a simple cut on the lost energy: $p_1^0 - p_2^0 \geq \Delta E$ where Δ is a dimensionless parameter, $\Delta \ll 1$ and $\Delta E \gg m_e$.

The typical magnitude of the $O(\alpha)$ corrections to the differential cross section is defined by three major factors:

$$\delta^{(1)} = \frac{d\sigma^{(1)}}{d\sigma^{(0)}} \sim \frac{\alpha}{2\pi} \cdot \ln\left(\frac{Q^2}{m_e^2}\right) \cdot \ln\Delta. \tag{7}$$

The enhancement by the so called large logarithm $L \equiv \ln\left(Q^2/m_e^2\right)$ and by the logarithm of the cut-off parameter make the size of the one-loop correction to be of the order of a few percent. Since the experimental uncertainties are well below this order, the one-loop corrections were taken into account in the data analysis [10].

Below we (re-)consider the following higher order QED effects:
1) higher order effects in vacuum polarization;
2) cut-off dependence of the photonic corrections;
3) light pair corrections in the leading logarithmic approximation;
4) complete next-to-leading $O\left(\alpha^2 L^1\right)$ corrections to the lepton line.

Running of the QED coupling constant can be naturally represented as

$$\alpha(Q^2) = \frac{\alpha(0)}{1 - \Pi(Q^2)}, \qquad \Pi(Q^2) = \Pi_e(Q^2) + \Pi_\mu(Q^2) + \Pi_{\rm hadr}(Q^2) + \dots \tag{8}$$

where $\alpha(0) \equiv \alpha \approx 1/137.036$. The magnitude of $\Pi(Q^2)$ for the range of momentum transfer under consideration is about 0.01. The bulk of the vacuum polarization effect comes from the one-loop e^+e^- pair insertion into the photon propagator,

$$\Pi_e(Q^2) = \frac{\alpha(0)}{\pi}\Big(\frac{1}{3}L - \frac{5}{9}\Big) + \Big(\frac{\alpha(0)}{\pi}\Big)^2\Big(\frac{1}{4}L + \zeta(3) - \frac{5}{24}\Big) + O(\alpha^3). \tag{9}$$

One can note that the $O\left(\alpha^2\right)$ contribution is of the next-to-leading order, since it contains only the first power of the large logarithm L. So it makes only a $\sim 10^{-5}$ effect well below the precision tag. The resummation of the vacuum polarization effect gives

$$\delta\sigma_{\rm vac.pol.} = \sigma^{(0)}\left(\frac{\alpha(Q^2)}{\alpha(0)}\right)^2 = \frac{\sigma^{(0)}}{|1 - \Pi(Q^2)|^2}. \tag{10}$$

Polarization of vacuum by virtual $\mu^+\mu^-$ pairs is not as large as by the e^+e^- ones. But in the bulk of the kinematical domain the suppression is only logarithmic. So,

$$\Pi_\mu(Q^2) = \frac{\alpha}{\pi}\left[\frac{v_\mu}{2}\left(1 - \frac{v_\mu^2}{3}\right)\ln\frac{v_\mu + 1}{v_\mu - 1} + \frac{v_\mu^2}{3} - \frac{8}{9}\right] + O\left(\alpha^2\right), \quad v_\mu \equiv \sqrt{1 + \frac{4m_\mu^2}{Q^2}} \tag{11}$$

has to be taken into account at least in the first order in α. For $Q^2 = 1$ GeV it reaches $2 \cdot 10^{-3}$.

Instead of the resummed geometrical series of Eq. 10, the A1 collaboration in Ref. [10] used exponentiation of the effect of the vacuum polarization by leptons, which is close numerically for the given Q^2 range.

Fig. 1 shows different contributions to the vacuum polarization correction $\delta_{\rm vac.pol.} = \delta\sigma_{\rm vac.pol.}/\sigma^{(0)}$. This figure was obtained with the help of the Fortran package `alphaQED` by F. Jegerlehner [24]. One can see that vacuum polarization by muons and hadrons contributes by up to one percent. That is a very large effect for the given precision tag. Moreover, the momentum dependence of the total vacuum polarization correction is different from the pure electron one.

As concerning the hadronic contribution to vacuum polarization, it can be either treated as a part of radiative corrections or as a part of the proton form factor. To our mind, the former treatment has two advantages. First, this contribution is always there as for point-like as well as for non-point-like particles. Second, in higher order corrections it is not factorized out as can be seen already in Eq. (10). From the first glance the hadronic contribution should not affect the value of the proton charge radius since it is defined at the zero momentum transfer, where this effect is vanishing. Nevertheless, the effect has a pronounced Q^2 dependence in the explored domain and it certainly affects the extrapolation to the zero momentum transfer point. For this reason we recommend to treat the hadronic vacuum polarization as a part of radiative corrections along this the corresponding leptonic contributions.

The Yennie-Frautschi-Suura theorem [14] proves that emission of each soft photon can be treated as an independent process. As the result, multiple emission of soft photons can be resummed into an exponent. By construction in this case, the maximal energy of each photon is limited independently. But in the given experimental set-up, we have a cut-off on the total lost energy. The corresponding effect was considered *e.g.* in Ref. [15]. For double soft photon emission in gives the following shift:

$$\mathrm{e}^{\delta_{\rm soft}} \to \mathrm{e}^{\delta_{\rm soft}} - \left(\frac{\alpha}{\pi}\right)^2\frac{\pi^2}{3}(L-1)^2\,. \tag{12}$$

At $Q^2 = 1$ GeV^2 this leads to a visible relative shift of the cross section of about $-3.5 \cdot 10^{-3}$. The same concerns emission of three and more soft photons, where the numerical effect of the proper cut on the lost energy is also not negligible. The proper exponentiation of radiative corrections in the leading logarithmic approximation is based on the exact solution of the renormalization group equation, see [16]. But for the practical application under consideration it is sufficient to compute effect order by order [17].

The contribution of e^+e^- pairs can be easily estimated with the help of the leading logarithmic approximation (LLA) in QED [16–19]:

$$\delta^{LLA}_{\rm pair} = \frac{2}{3}\left(\frac{\alpha}{2\pi}L\right)^2 P^{(0)}_\Delta + \frac{4}{3}\left(\frac{\alpha}{2\pi}L\right)^3 \left\{\left(P^{(0)} \otimes P^{(0)}\right)_\Delta + \frac{2}{9}P^{(0)}_\Delta\right\} + O\left(\alpha^2 L, \alpha^4 L^4\right) \tag{13}$$

where the so-called Δ-parts of splitting functions (see *e.g.* Refs. [17, 18]) read

$$P^{(0)}_\Delta = 2\ln\Delta + \frac{3}{2}, \qquad \left(P^{(0)} \otimes P^{(0)}\right)_\Delta = \left(P^{(0)}_\Delta\right)^2 - \frac{\pi^2}{3}. \tag{14}$$

Note that in the third order in α we have an effect due to simultaneous (either virtual or soft) radiation of a pair and a photon. To have a better control on the precision level, we can include also the next-to-leading pair corrections in the order $O\left(\alpha^2 L\right)$ where some enhancement due to the experimental cut-off takes place. The corresponding effect will be estimated below.

In order to control the precision of theoretical estimates we can compute the complete set of next-to-leading order (NLO) corrections to the given process by means of the renormalization group approach. The NLO QED structure functions were first introduced in [20]. The corresponding fragmentation functions were used in Refs. [21, 22] to evaluate NLO corrections to the muon decay spectrum. Here we can follow the approach developed in Ref. [23], where NLO QED corrections were computed in a similar set-up for Bhabha scattering.

The relevant photonic and e^+e^- pair contributions to the NLO electron structure (str) and fragmentation (frg) functions have the form[1]

$$\begin{aligned} D^{\rm str,frg}_{ee}(z) &= \delta(1-z) + \frac{\alpha}{2\pi}\Big(LP^{(0)}(z) + d_1(z)\Big) + \left(\frac{\alpha}{2\pi}\right)^2\Big(\frac{1}{2}L^2 P^{(0)} \otimes P^{(0)}(z) + \frac{1}{3}L^2 P^{(0)}(z) \\ &+ LP^{(0)} \otimes d_1(z) + LP^{(1,\gamma)\rm str,frg}_{ee}(z) + LP^{(1,\rm pair)str,frg}_{ee}(z)\Big) + O\left(\alpha^2 L^0, \alpha^3\right). \end{aligned} \tag{15}$$

Explicit expressions for splitting functions $P^{(n)}_{ee}$ and d_1 can be found in [23]. The master formula for NLO photonic corrections to elastic electron-proton scattering reads

$$d\sigma = \int_{\bar z}^{1} dz D^{\rm str}_{ee}(z)\Big(d\sigma^{(0)}(z) + d\bar\sigma^{(1)}(z) + O\left(\alpha^2 L^0\right)\Big)\int_{\bar y}^{1}\frac{dy}{Y}D^{\rm frg}_{ee}\left(\frac{y}{Y}\right), \tag{16}$$

where $d\bar\sigma^{(1)}$ is the $O(\alpha)$ correction to the ep scattering with a "massless electron", calculated using the $\overline{\rm MS}$ scheme to subtract the lepton mass singularities. The energy fraction of the incoming parton is z, and Y is the the energy fraction of the outgoing (observed) electron. As concerning the factorization scale, it is natural to choose it to be equal to the momentum transfer: $L \equiv \ln(Q^2/m_e^2)$.

Here we are interested in the contributions due to virtual and soft photons, so both integrals have the same lower limit being equal to $1 - \Delta$. First we can perform convolution of the structure and fragmentation functions entering Eq. (16) with each other $D^{\rm str}_{ee} \otimes D^{\rm frg}_{ee}(z)$. If $z = 1 - \Delta$ and $\Delta \ll 1$, the result of the convolution gives the probability density to find such a situation where one looses in total due to photon emission $\Delta E_{\rm beam}$ from the total energy of the process under consideration.

[1] We dropped the singlet channel contributions which are suppressed in the given experimental set-up.

Convolution of the function found above with the Born part of the kernel cross section gives us the corresponding cross section (with the upper limit on the lost energy):

$$d\sigma^{\rm NLO} = d\sigma^{(0)}(1)\Big\{1 + 2\frac{\alpha}{2\pi}\Big[LP^{(0)}_{\Delta} + (d_1)_{\Delta}\Big] + 2\Big(\frac{\alpha}{2\pi}\Big)^2\Big[L^2\left(P^{(0)}\otimes P^{(0)}\right)_{\Delta} + \frac{1}{3}L^2P^{(0)}_{\Delta}$$
$$+2L(P^{(0)}\otimes d_1)_{\Delta} + L(P^{(1,\gamma)}_{ee})_{\Delta} + L(P^{(1,{\rm pair})}_{ee})_{\Delta}\Big]\Big\} + d\bar{\sigma}^{(1)}(1)\,2\frac{\alpha}{2\pi}LP^{(0)}_{\Delta} + O\left(\alpha^3L^3\right), \quad (17)$$

$$(d_1)_{\Delta} = -2\ln^2\Delta - 2\ln\Delta + 2,$$
$$(P^{(0)}\otimes d_1)_{\Delta} = -4\ln^3\Delta - 7\ln^2\Delta + \ln\Delta\,(1+8\zeta(2)) + 3 - 8\zeta(3) + 4\zeta(2),$$
$$(P^{(1,\gamma)}_{ee})_{\Delta} = \frac{3}{8} - 3\zeta(2) + 6\zeta(3), \qquad (P^{(1,{\rm pair})}_{ee})_{\Delta} = -\frac{20}{9}\ln\Delta - \frac{1}{6} - \frac{4}{3}\zeta(2). \quad (18)$$

The values of the Riemann zeta function are $\zeta(2) = \pi^2/6$ and $\zeta(3) \approx 1.202$. The Δ-parts of the structure and fragmentation splitting functions $(P^{(1,\gamma({\rm pair})){\rm str}}_{ee})_{\Delta}$ and $(P^{(1,\gamma({\rm pair})){\rm frg}}_{ee})_{\Delta}$ coincide, so the notation is simplified. By construction in the $\overline{\rm MS}$ scheme, the complete first order correction is reproduced since

$$d\bar{\sigma}^{(1)}(1) \;=\; d\sigma^{(1)}(1) - 2d\sigma^{(0)}(1)\frac{\alpha}{2\pi}\Big[LP^{(0)}_{\Delta} + (d_1)_{\Delta}\Big]. \quad (19)$$

The factor 2 before the subtracted term on the right hand side reflects the presence of mass singularities in both the initial and final state corrections.

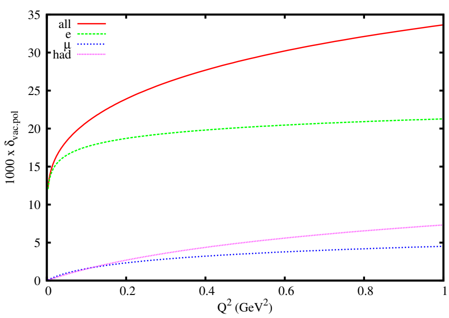

Figure 1. Vacuum polarization corrections due to electrons (e), muons (μ), hadrons (had), and the combined effect (all).

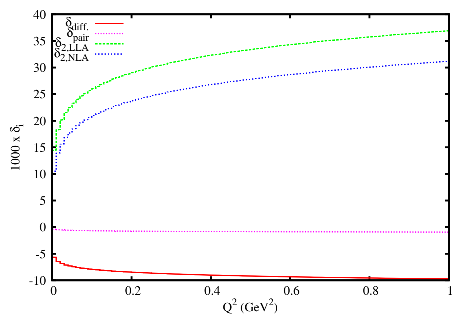

Figure 2. Relative higher order QED corrections to electron line in *ep* scattering cross section *vs* momentum transferred squared.

Relative QED corrections to the electron line

$$\delta_i = \frac{d\sigma^{(i)}}{d\sigma^{(0)}} \quad (20)$$

are presented in Fig. 2. Index i runs over:
a) “2,LLA”, *i.e.* pure photonic $O\left(\alpha^2L^2\right)$ corrections from Eq. (17),
b) “2,NLA”, *i.e.* the sum of pure photonic $O\left(\alpha^2L^2\right)$ and $O\left(\alpha^2L^1\right)$ corrections from Eq. (17),
c) “pair”, *i.e.* the leading log pair corrections from Eq. (13) supplemented by subleading pair corrections extracted from Eq. (17),
d) “diff.”, *i.e.* the shift from the exponentiated one-loop result:

$$\delta_{\rm diff.} \;=\; \frac{d\sigma^{\rm NLO}}{d\sigma^{(0)}(1)} + \delta^{(3)}_{\rm LLA} + \delta^{(3)}_{\rm LLA,pair} + \delta^{(4)}_{\rm LLA} - \exp\{\delta^{(1)}\}. \quad (21)$$

We claim that the advanced treatment of RC to elastics electron-proton scattering, which is presented above, should be applied in the analysis of experimental data in future experiments. An adequate treatment of all other relevant effects (double photon exchange, radiative corrections to the proton line, details of the experimental set-up, *etc.*) is also required.

4 Conclusions

Two types of effects due to radiative corrections are discussed. We have shown that these effects might be relevant for modern high-precision experiments.

We are grateful to S. Karshenboim for fruitful discussions.

References

[1] A.B. Arbuzov and T.V. Kopylova, Phys. Part. Nucl. Lett. **11**, 339 (2014).
[2] A.B. Arbuzov and T.V. Kopylova, Eur. Phys. J. C **75**, 603 (2015).
[3] K.A. Olive *et al.*, Particle Data Group Collaboration, Chin. Phys. C **38**, 090001 (2014).
[4] S. Actis *et al.*, Working Group on Radiative Corrections and Monte Carlo Generators for Low Energies Collaboration, Eur. Phys. J. C **66**, 585 (2010).
[5] B.L. Roberts, Chin. Phys. C **34**, 741 (2010).
[6] T. Mibe, J-PARC g-2 Collaboration, Nucl. Phys. Proc. Suppl. **218**, 242 (2011).
[7] I.F. Ginzburg, G.L. Kotkin, S.I. Polityko and V.G. Serbo, Phys. Lett. B **286**, 392 (1992).
[8] A.I. Akhiezer and V.B. Berestetskii, *Elements of Quantum Electrodynamics* (Oldbourne Press, London, 1962).
[9] J. S. Schwinger, Phys. Rev. **73**, 416 (1948).
[10] J. C. Bernauer *et al.* [A1 Collaboration], Phys. Rev. C **90**, 015206 (2014).
[11] A. Antognini *et al.*, Science **339**, 417 (2013).
[12] J. Arrington, P. G. Blunden and W. Melnitchouk, Prog. Part. Nucl. Phys. **66**, 782 (2011).
[13] G. Lee, J.R. Arrington and R.J. Hill, Phys. Rev. D **92**, 013013 (2015).
[14] D.R. Yennie, S.C. Frautschi and H. Suura, Annals Phys. **13**, 379 (1961).
[15] A.B. Arbuzov, E.A. Kuraev and B.G. Shaikhatdenov, Mod. Phys. Lett. A **13**, 2305 (1998).
[16] E.A. Kuraev and V.S. Fadin, Sov. J. Nucl. Phys. **41**, 466 (1985).
[17] A.B. Arbuzov, Phys. Lett. B **470**, 252 (1999).
[18] M. Skrzypek, Acta Phys. Polon. B **23**, 135 (1992).
[19] A. B. Arbuzov, JHEP **0107**, 043 (2001).
[20] F.A. Berends, W.L. van Neerven and G.J. H. Burgers, Nucl. Phys. B **297** 429 (1988), *Erratum-ibid.* B **304**, 921 (1988).
[21] A. Arbuzov and K. Melnikov, Phys. Rev. D **66**, 093003 (2002).
[22] A. Arbuzov, JHEP **0303**, 063 (2003).
[23] A.B. Arbuzov and E.S. Scherbakova, JETP Lett. **83**, 427 (2006).
[24] F. Jegerlehner, Nuovo Cim. C **034S1**, 31 (2011).