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TWO–PHOTON REACTIONS BEYOND ONE–LOOP

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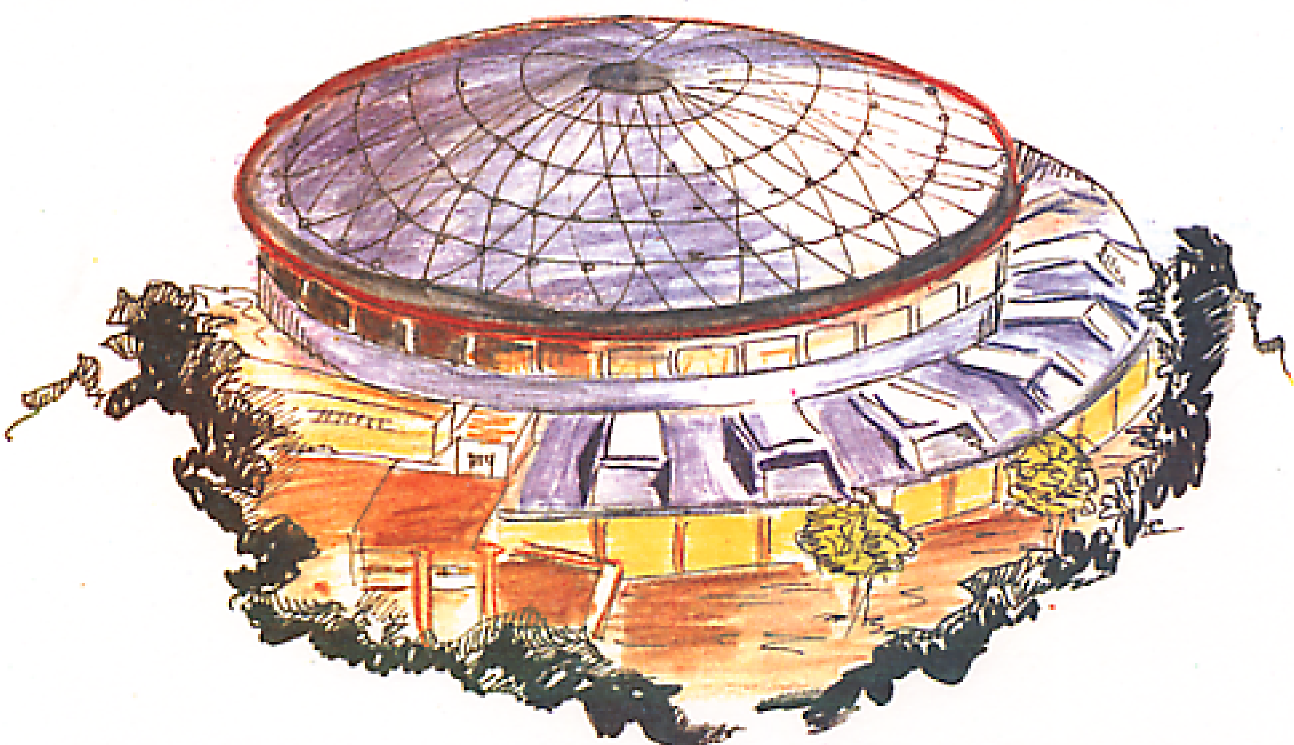

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TWO–PHOTON REACTIONS BEYOND ONE–LOOP

S. Bellucci[a]

INFN-Laboratori Nazionali di Frascati, P.O.Box 13, 00044 Frascati, Italy.

Abstract

We review the recent progress in the calculation of the amplitude for $\gamma\gamma \rightarrow \pi^0\pi^0$ to two loops in chiral perturbation theory. We match the low-energy amplitude in chiral perturbation theory with the result of the dispersion theoretic analysis. The neutral pion polarizabilities are also given to two-loop accuracy. Then, the results are compared with the dispersion relation calculation of the pion polarizabilities.

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a) e-mail: Bellucci@lnf.infn.it

The process producing a pion pair in the fusion of two photons has received a lot of attention in the recent years. It provides a very important test of chiral perturbation theory (CHPT) [1]-[5]. In this sense, the production of neutral pion pairs is the most interesting channel. This is because the Born amplitude vanishes in this case. Hence the one-loop scattering amplitude is finite and does not depend on the free parameters of the chiral lagrangian [6, 7]. It is by now a well-known fact that the one-loop cross section for $\gamma\gamma \to \pi^0\pi^0$ in CHPT [6, 7] does not agree with the experimental measurements at Crystal Ball [8], as well as with calculations based on dispersion relations [9]-[15], even at low-energy. We calculated recently the low-energy $\gamma\gamma \to \pi^0\pi^0$ amplitude to two-loops in CHPT [16] and obtained a prediction that agrees with the Crystal Ball data. Also, the low-energy CHPT amplitude compares very well with the dispersive analysis of $\gamma\gamma \to \pi^0\pi^0$ by Donoghue and Holstein [12].

The sum of the electric and magnetic polarizabilities of the neutral pion vanishes to lowest order in CHPT [17]. The value of this sum was estimated also by a sum rule, and it turned out to be different from zero [18]. We showed some time ago [19] that a vector dominance model preserving the chiral symmetry of QCD at low energy yields a value close to the one obtained from the sum rule. We obtained recently a refinement of this prediction, by calculating the two-loop Compton scattering amplitude in CHPT [16]. The largest modification in the polarizabilities, with respect to the one-loop order value, is given by the omega resonance exchange which accounts for a large fraction of the neutral pion sum rule [19], and the contribution from the chiral logarithms is small [16].

The $\gamma\gamma \to \pi^0\pi^0$ cross section for off-shell photons has been calculated recently [20] in the framework of CHPT. There, it has been shown that the measurement of the azymuthal correlations in the process $e^+e^- \to e^+e^-\pi^0\pi^0$ allows to test the higher order CHPT corrections independently from the measurement of the cross-section.

Recently, CHPT has been reformulated to include [21] into each order additional terms which in the standard CHPT are of higher order. Within this generalization of the chiral expansion of the amplitude, the process $\gamma\gamma \to \pi^0\pi^0$ has been analyzed [22].

Gauge symmetry and Lorentz invariance can be used to write the scattering matrix element

$$< \pi^0(p_1)\pi^0(p_2)\text{out} \mid \gamma(q_1)\gamma(q_2)\text{in} >= i(2\pi)^4\delta^4(P_f - P_i)T^N \quad , \qquad (0.1)$$

with

$$
\begin{aligned}
T^N &= e^2 \epsilon_1^\mu \epsilon_2^\nu V_{\mu\nu} \ , \\
V_{\mu\nu} &= i \int dx e^{-i(q_1 x + q_2 y)} < \pi^0(p_1)\pi^0(p_2)\text{out} \mid T j_\mu(x) j_\nu(y) \mid 0 >,
\end{aligned}
\tag{0.2}
$$

where j_μ is the electromagnetic current, and $\alpha = e^2/4\pi \simeq 1/137$, as follows:

$$
\begin{aligned}
V_{\mu\nu} &= A(s,t,u)T_{1\mu\nu} + B(s,t,u)T_{2\mu\nu} + \ , \\
T_{1\mu\nu} &= \frac{s}{2} g_{\mu\nu} - q_{1\nu} q_{2\mu} \ , \\
T_{2\mu\nu} &= 2s\Delta_\mu \Delta_\nu - \nu^2 g_{\mu\nu} - 2\nu(q_{1\nu}\Delta_\mu - q_{2\mu}\Delta_\nu) \ , \\
\Delta_\mu &= (p_1 - p_2)_\mu \ ,
\end{aligned}
\tag{0.3}
$$

in terms of the standard Mandelstam variables

$$
\begin{aligned}
s &= (q_1 + q_2)^2, \ t = (p_1 - q_1)^2, \ u = (p_2 - q_1)^2 \ , \\
\nu &= t - u.
\end{aligned}
\tag{0.4}
$$

One can go from the analytic functions A and B of the variables s, t and u, symmetric under crossing $(t, u) \to (u, t)$, to the helicity amplitudes in the following way:

$$
\begin{aligned}
H_{++} &= A + 2(4M_\pi^2 - s)B \ , \\
H_{+-} &= \frac{8(M_\pi^4 - tu)}{s} B \ .
\end{aligned}
\tag{0.5}
$$

In Ref. [16] the renormalization procedure is formulated in the minimal subtraction scheme, and the expressions of the order E^6 renormalized amplitudes involve three parameters, i.e. a_1^r, a_2^r and b^r

$$
A_6 = \frac{a_1^r M^2 + a_2^r s}{(16\pi^2 F^2)^2} + \cdots \ , \tag{0.6}
$$

$$
B_6 = \frac{b^r}{(16\pi^2 F^2)^2} + \cdots \ , \tag{0.7}
$$

where the ellipses stand for finite contributions from the loop-integrals. Here F is the pion decay constant in the chiral limit, $F_\pi = F(1 + O(\hat{m})), F_\pi \simeq 93$ MeV, and the physical pion mass is

$$
\begin{aligned}
M_\pi^2 &= M^2(1 + O(\hat{m})) \ , \\
M^2 &= 2\hat{m}B,
\end{aligned}
\tag{0.8}
$$

Table 1: Resonance contributions to the coupling constants a_1^r, a_2^r and b^r. Column 6 contains the sums of those contributions which have a definite sign.

	I^R					I^R	
I^r	ω	ρ^0	ϕ	$A(1^{+-})$	$\sum_R I^R$	$S(0^{++})$	f_2
a_1^r	-33.2	-6.1	-0.1	0.0	-39	± 0.8	∓ 4.1
a_2^r	12.5	2.3	$\simeq 0$	-1.3	13	± 1.3	± 1.0
b^r	2.1	0.4	$\simeq 0$	0.7	3	0.0	± 0.5

where B is related to the order parameter $< 0 \mid \bar{q}q \mid 0 >$, and in the isospin symmetry limit $m_u = m_d = \hat{m}$.

The three unknown parameters a_1^r, a_2^r and b^r are estimated by resonance saturation [16] and the results are displayed in table 1. The ω exchange yields the dominant contribution.

The low-energy constants $h_\pm^r$ and h_s^r corresponding to the helicity amplitudes H_{++} and H_{+-} read

$$\begin{aligned} H_{++}^{\rm 2loops} &= \frac{1}{(16\pi^2F^2)^2}\left\{h_+^r M^2 + h_s^r\, s\right\} + \cdots \ , \\ H_{+-}^{\rm 2loops} &= \frac{8(M^4 - tu)}{s(16\pi^2F^2)^2} h_-^r + \cdots \ , \\ h_+^r &= a_1^r + 8b^r \ , \ h_s^r = a_2^r - 2b^r \ , \ h_-^r = b^r \ . \end{aligned} \tag{0.9}$$

From column 6 in table 1 the central values of these couplings are obtained in Ref. [16], where a 30% uncertainty is associated to the contributions generated by the vector and axial-vector exchange, and a 100% error to the contributions from scalars and from f_2. Adding these errors in quadrature, one finds [16]

$$\begin{aligned} h_+^r(M_\rho) &= -14 \pm 5 \ , \\ h_s^r(M_\rho) &= 7 \pm 3 \ , \\ h_-^r(M_\rho) &= 3 \pm 1 \ . \end{aligned} \tag{0.10}$$

The values of h_+^r and h_s^r are not affected by the tensor exchange, since the corresponding coupling is purely D-wave. Scalars do not contribute to the value of h_-^r.

In Ref. [23], these couplings have been determined i) from vector-meson exchange and using nonet-symmetry, and ii) from the chiral quark model, with the result

$$(h_+^r, h_s^r, h_-^r)|_{\mu=M_\rho} = \begin{cases} (-18, 9, 2) & \text{vector-mesons (nonet)} \\ (-12, 6, 2) & \text{chiral quark model} \end{cases} \quad (0.11)$$

which agrees within the uncertainties with the values in (0.10).

The following analytic results are obtained in Ref. [16] for the amplitude A to two loops:

$$A = \frac{4\bar{G}_\pi(s)}{sF_\pi^2}(s - M_\pi^2) + U_A + P_A + O(E^4). \quad (0.12)$$

The unitary part U_A contains s, t and u-channel cuts, and P_A is a linear polynomial in s. Explicitly,

$$\begin{aligned} U_A &= \frac{2}{sF_\pi^4}\bar{G}(s)\left[(s^2 - M_\pi^4)\bar{J}(s) + C(s, \bar{l}_i)\right] + \frac{\bar{l}_\Delta}{24\pi^2 F_\pi^4}(s - M_\pi^2)\bar{J}(s) \\ &+ \frac{(\bar{l}_2 - 5/6)}{144\pi^2 sF_\pi^4}(s - 4M_\pi^2)\left\{\bar{H}(s) + 4\left[s\bar{G}(s) + 2M_\pi^2(\bar{\bar{G}}(s) - 3\,\bar{\bar{J}}(s))\right]d_{00}^2\right\} \\ &+ \Delta_A(s, t, u) \ , \end{aligned} \quad (0.13)$$

with

$$\begin{aligned} C(s, \bar{l}_i) &= \frac{1}{48\pi^2}\left\{2(\bar{l}_1 - 4/3)(s - 2M_\pi^2)^2 + (\bar{l}_2 - 5/6)(4s^2 - 8sM_\pi^2 + 16M_\pi^4)/3 \right. \\ &\left. -3M_\pi^4\bar{l}_3 + 12M_\pi^2(s - M_\pi^2)\bar{l}_4 - 12sM_\pi^2 + 15M_\pi^4\right\} \ , \\ d_{00}^2 &= \frac{1}{2}(3\cos\theta^2 - 1) \, . \end{aligned} \quad (0.14)$$

The loop-functions $\bar{J}$ etc. are given in appendix C of Ref. [16].

The polynomial part is

$$\begin{aligned} P_A &= \frac{1}{(16\pi^2 F_\pi^2)^2}[a_1 M_\pi^2 + a_2 s] \ , \\ a_1 &= a_1^r + \frac{1}{18}\left\{4l^2 + l(8\bar{l}_2 + 12\bar{l}_\Delta - \frac{4}{3}) - \frac{20}{3}\bar{l}_2 + 12\bar{l}_\Delta + \frac{110}{9}\right\} \ , \\ a_2 &= a_2^r - \frac{1}{18}\left\{l^2 + l(2\bar{l}_2 + 12\bar{l}_\Delta + \frac{2}{3}) - \frac{5}{3}\bar{l}_2 + 12\bar{l}_\Delta + \frac{697}{144}\right\} \ , \\ l &= \ln\frac{M^2}{\mu^2} \ . \end{aligned} \quad (0.15)$$

Here μ denotes the renormalization scale, the $\bar{l}_i$ are the renormalized coupling constants of the $O(E^4)$ lagrangian [2], and $\bar{l}_\Delta = \bar{l}_6 - \bar{l}_5$. The values of these coupling constants can be found in column 2 of table 1 of Ref. [16]. The result for B is

$$B = U_B + P_B + O(E^2) \ , \tag{0.16}$$

with unitary part

$$U_B = \frac{(\bar{l}_2 - 5/6)\bar{H}(s)}{288\pi^2 F_\pi^4 s} + \Delta_B(s,t,u) \ . \tag{0.17}$$

The polynomial is obtained in Ref. [16]

$$\begin{aligned} P_B &= \frac{b}{(16\pi^2 F_\pi^2)^2} \ , \\ b &= b^r - \frac{1}{36}\left[l^2 + l(2\bar{l}_2 + \frac{2}{3}) - \frac{\bar{l}_2}{3} + \frac{393}{144}\right] \ . \end{aligned} \tag{0.18}$$

The integrals $\Delta_{A,B}(s,t,u)$ contain contributions that very small for the cross sections below $\sqrt{s} \leq 400$ MeV, both for $\gamma\gamma \to \pi^0\pi^0$ (0.1% at 400 MeV) and for the crossed channel $\gamma\pi^0 \to \gamma\pi^0$ (1.5% at 400 MeV) [16].

The cross section $\gamma\gamma \to \pi^0\pi^0$ receives a substantial correction near threshold due to $\pi\pi$ final-state interactions – which are absent in Compton scattering. Yet, it is shown in Ref. [16] that the two-loop contributions are not small in this channel. Since in the one-loop approximation the amplitude H_{++} is one order of magnitude larger in the $\gamma\gamma \to \pi^0\pi^0$ channel than at Compton threshold, even very small corrections in $\gamma\gamma \to \pi^0\pi^0$ may appear large in Compton scattering [19]. The result of the two-loop calculation and the one-loop approximation differ by one order of magnitude already near threshold [16]. This is mainly due to the effect of the low-energy constant h_-^r in H_{+-} (omega-exchange in the language of resonance saturation [19]).

The low-energy limit of the coupling with the photon in the Compton amplitude for a composite system is characterized (among other parameters) by the electric and magnetic polarizabilities. One can test the hadron dynamics through experiments on the hadron polarizabilities [18]. The expansion of the amplitude for charged pion Compton scattering,

$$\gamma(q_1)\pi^+(p_1) \to \gamma(q_2)\pi^+(p_2) \ , \tag{0.19}$$

near threshold reads

$$T^C = 2\left[\vec{\epsilon}_1 \cdot \vec{\epsilon}_2^{\,\star}\left(\frac{\alpha}{M_\pi} - \bar{\alpha}_\pi \omega_1 \omega_2\right) - \bar{\beta}_\pi (\vec{q}_1 \times \vec{\epsilon}_1) \cdot (\vec{q}_2 \times \vec{\epsilon}_2^{\,\star}) + \cdots\right] \tag{0.20}$$

with $q_i^\mu = (\omega_i, \vec{q}_i)$. For neutral pions, one has, in terms of A and B,

$$\begin{aligned} \bar{\alpha}_{\pi^0} &= \frac{\alpha}{2M_\pi}(A + 16M_\pi^2 B)|_{s=0,t=M_\pi^2} \ , \\ \bar{\beta}_{\pi^0} &= -\frac{\alpha}{2M_\pi} A|_{s=0,t=M_\pi^2} . \end{aligned} \tag{0.21}$$

Below we denote

$$(\alpha \pm \beta)^C = \bar{\alpha}_\pi \pm \bar{\beta}_\pi \ ,$$
$$(\alpha \pm \beta)^N = \bar{\alpha}_{\pi^0} \pm \bar{\beta}_{\pi^0} \ . \qquad (0.22)$$

The pion polarizabilities have been estimated[1] through dispersion sum rules [18]

$$(\alpha + \beta)^C = 0.39 \pm 0.04 \ ,$$
$$(\alpha - \beta)^C = 10 \pm 3 \ ,$$
$$(\alpha + \beta)^N = 1.04 \pm 0.07 \ ,$$
$$(\alpha - \beta)^N = -10 \pm 4 \ . \qquad (0.23)$$

The charged pion polarizabilities have been determined in an experiment on the radiative pion-nucleus scattering $\pi^- A \to \pi^- \gamma A$ [24] and in the pion photoproduction process $\gamma p \to \gamma \pi^+ n$ [25]. Assuming the constraint $(\alpha+\beta)^C = 0$ the two experiments yield

$$(\alpha - \beta)^C = \begin{cases} 13.6 \pm 2.8 & [24] \\ 40 \pm 24 & [25] \ . \end{cases} \qquad (0.24)$$

Relaxing the constraint $(\alpha + \beta)^C = 0$, one obtains from the Serpukhov data

$$(\alpha + \beta)^C = 1.4 \pm 3.1(\text{stat.}) \pm 2.5(\text{sys.}) \ [26] \ ,$$
$$(\alpha - \beta)^C = 15.6 \pm 6.4(\text{stat.}) \pm 4.4(\text{sys.}) \ [26] \ . \qquad (0.25)$$

At one-loop one has [17, 27]

$$\bar{\alpha}_{\pi^0} = -\bar{\beta}_{\pi^0} = -\frac{\alpha}{96\pi^2 M_\pi F^2} = -0.50 \ . \qquad (0.26)$$

At order $O(E^6)$ we found [16]

$$\bar{\alpha}_{\pi^0} = -0.50 + 0.21 - 0.07 \simeq -0.35 \ ,$$
$$\bar{\beta}_{\pi^0} = 0.50 + 0.79 + 0.24 \simeq 1.50 \qquad (0.27)$$

where the three contributions that add up to the final results on the r.h.s. are the one-loop, the resonance, and the two-loop contributions, respectively. There is a

[1]The values of the polarizabilities are in units of 10^{-4}fm^3 in what follows.

large contribution from the resonance exchange. Our results saturates the forward sum rule $(\alpha+\beta)^N = 1.04 \pm 0.07$ in (0.23)

$$(\alpha+\beta)^N = 1.0 + 0.16 \simeq 1.15 \quad . \tag{0.28}$$

Information on the charged pion polarizabilities may be obtained from $\gamma\gamma \to \pi^+\pi^-$ data [27]. The low-energy constant $\bar{l}_\Delta$ appears as the only free parameter order in the $O(E^4)$ amplitude, as well as in the leading-order expression for $\bar{\alpha}_\pi$ and $\bar{\beta}_\pi$. A fit to the cross section then determines $\bar{\alpha}_\pi$ and $\bar{\beta}_\pi$. The result [27] $\bar{l}_\Delta = 2.3 \pm 1.7$ corresponds to numerical value for the leading-order $\bar{\alpha}_\pi = 2.7 \pm 0.4$, plus systematic uncertainties due to the $O(E^6)$ corrections. For the charged pions, a two-loop calculation is not yet available. The charged pion polarizabilities are given beyond the one-loop order by including the meson resonance contribution in Refs. [28, 19]

The construction of unitarized S-wave amplitudes for $\gamma\gamma \to \pi\pi$ which contain $(\alpha-\beta)^{C,N}$ as adjustable parameters has been carried out in Ref. [15]. In this case, only $(\alpha-\beta)^{C,N}$ can be determined from the data [8, 29], with the result

$$\begin{aligned} (\alpha-\beta)^C &= \quad 4.8 \pm 1.0 \;\; [15] \;, \\ (\alpha-\beta)^N &= -1.1 \pm 1.7 \;\; [15] \,. \end{aligned} \tag{0.29}$$

The value (0.29) for $(\alpha-\beta)^N$ is consistent with the two-loop result for the neutral pion, whereas the corresponding calculation for charged pions is not available and so it cannot be compared with the value (0.29) for $(\alpha-\beta)^C$.

Finally, it is interesting to compare the chiral expansion [16] with the dispersive calculation carried out by Donoghue and Holstein [12]. The two representations of the S-wave amplitude agree numerically very well below $E = 0.4\ GeV$. In the dispersive method, higher order terms are partially summed up. The agreement indicates that yet higher orders in the chiral expansion do not affect much the threshold amplitude.

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