\( \pi^+ \pi^- \) ATOM IN CHIRAL THEORIES

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ABSTRACT

We review the existing theoretical approaches to the study of the observables of hadronic atoms. In the relativistic perturbative approach based on the Bethe-Salpeter equation and Chiral Perturbation Theory we derive the general expression for the \( \pi^+ \pi^- \) atom lifetime. The lowest-order correction to the relativistic Deser-type formula for the atom width is predicted to be \((6.1 \pm 3.1)\%\).

The study of the pion-pion scattering which forms one of the basic building blocks in the hierarchy of strong interaction processes, enables one to gain a deeper insight in the nature of strong interactions. According to common belief, low-energy interactions of pions are described within the Chiral Perturbation Theory (ChPT) \(^1\) which exploits a full content of global QCD symmetries. In the standard scheme with a "large" condensate, the \( S \)-wave \( \pi\pi \) scattering lengths are predicted to be \( a_0^0 = 0.217 \) and \( a_0^0 - a_2^0 = 0.258 \) in units of the inverse charged pion mass \(^2\). The calculations within the Generalized ChPT with a small quark condensate which contain more parameters, lead to a most likely value \( a_0^0 = 0.27 \) \(^3\). Despite a significant difference between these numbers, both results for the scattering length \( a_0^0 \) are compatible with the experimental value \( a_0^0 = 0.26 \pm 0.05 \) \(^4\). Consequently, a precise measurement of \( \pi\pi \) scattering lengths will be an excellent test of the ChPT.

The existing experimental information on the \( S \)-wave \( \pi\pi \) scattering lengths is extracted either from the data on the process \( \pi N \to \pi\pi N \) or from the study of \( K_{e4} \) decay (for the review of the recent status of \( \pi\pi \) experiments see, e.g. \(^5\)). The forthcoming DIRAC experiment PS212 at CERN is aimed at the high-precision measurement of the lifetime of the \( \pi^+ \pi^- \) atom \(^6\). This will allow for the determination of the quantity \( a_0^0 - a_2^0 \) with an accuracy of 5\% and thus will provide a decisive probe of the predictions of ChPT. The experiments on the measurement of the hadronic atom observables are carried out also at PSI, KEK, Frascati, Uppsala.

The experimental study of \( \pi^+ \pi^- \) atom characteristics provides a direct information about the strong \( \pi\pi \) amplitudes near threshold since the average momenta \((\sim \alpha m_\pi)\) of constituents in the atom is much less than either of two typical scales characterizing strong interactions in chiral theories \((\sim m_\pi \text{ and } 4\pi F_\pi \sim 1 \text{ GeV})\). The huge difference between the bound-state and strong scales leads to a clear-cut factorization of strong and
Electromagnetic interactions in the observables of the pionium which is already explicit in the lowest-order Deser formula for the ground-state atom lifetime.

\[ \tau_1^{-1} = \frac{16\pi}{9} \left( \frac{2\Delta m_\pi}{m_\pi} \right)^{1/2} (a_0^0 - a_0^2)^2 |\Psi_1(0)|^2 \]

where \( \Psi_1(0) \) denotes the value of the nonrelativistic Coulomb ground-state wave function at the origin. Note that the Coulomb interaction which is mainly responsible for the atom formation, and strong interactions which lead to its decay, enter separately into the expression of the lifetime through the Coulomb wave function and the scattering lengths, respectively. However, in the view of the forthcoming DIRAC high-precision experiment it is important to evaluate corrections to this formula in order to provide a careful determination of the strong scattering lengths. In particular, one can readily observe from Eq. (1) that the change of the "reference mass" from the charged to the neutral pion mass in the definition of strong scattering lengths leads to a \( \sim 12\% \) variation in the lifetime. In order to avoid this ambiguity one has to perform a choice of the "reference mass" in the ideal, isotopically symmetric world and evaluate all mass shift corrections which emerge due to the \( m_{\pi^+} - m_{\pi^0} \) mass difference in the pionium observables. All other corrections of an electromagnetic or a strong origin should also be calculated.

Historically, the first attempts to calculate the corrections to the pionium lifetime within the field-theoretical approaches were focused on the evaluation of the quantities entering into Eq. (1): the scattering lengths and the wave function at the origin. The value of \( a_0^0 - a_0^2 \) was calculated in various low-energy phenomenological models. The \( \pi\pi \) strong interaction potential was constructed and the corrections to the pionium energy and the value of the wave function at the origin due to strong interactions, vacuum polarization, finite size of pion were evaluated, taking into account only the discrete Coulomb spectrum. The vacuum polarization effect in the pionium observables was studied in the quasipotential approach. Retardation correction to the pionium lifetime was calculated on the basis of Bethe-Salpeter (BS) approach. Electromagnetic radiative corrections to the \( \pi\pi \) scattering lengths and the pionium lifetime were evaluated. Note however that the consistent approach to the calculation of bound-state characteristics can not be confined solely to the quantities entering into the Deser formula. Instead, to this end a systematic field-theoretical framework is needed which is free of any double counting problems. In the several recent papers an attempt was made to carry out a comprehensive field theoretical analysis of the pionium problem. Namely, in the basis of the 3-dimensional constraint theory bound-state equations the mass shift, radiative and second-order strong corrections have been evaluated. In our previous papers a systematic perturbative approach based on the BS equation is constructed, and a complete set of the lowest-order corrections to the pionium lifetime is evaluated. The vacuum polarization effect on the pionium observables is evaluated within the nonrelativistic QED.

In a nonrelativistic potential approach to the problem the expression for the lifetime, including the corrections, is written in terms of physical scattering lengths in charged and neutral channels in the presence of Coulomb interactions. This is a counterpart of the lowest-order Deser-type factorization: the bound-state observables are expressed in terms of the quantities characterizing the scattering process, and these quantities differ from the purely strong ones by a small and calculable mass shift and electromagnetic corrections. In order to evaluate these corrections, one needs the \( \pi\pi \) potential to be given explicitly. Assuming local and energy-independent strong potentials in the isospin symmetry limit,
one reproduces the $\pi\pi$ scattering phase shifts calculated from ChPT. At the next step the Coulomb potential is added and the mass shift effect is included via assigning the physical masses to pions, the strong potential remaining unchanged. The sign of the largest mass shift correction in the scattering lengths and the pionium lifetime calculated in this manner \cite{19} turns out to be opposite to that emerging from the field-theoretical calculations \cite{13,14,17} provided the charged pion mass is chosen to be the ”reference mass”. Owing to the derivative character of pion couplings in the chiral Lagrangian, a possible reason for this discrepancy might be an explicit energy independence of strong potentials used in these calculations.

To summarize, one needs an accurate theoretical estimate of the lifetime of the pionium which will be measured in a high precision DIRAC experiment. The underlying dynamics of pions should be described by the Lagrangian of ChPT in order to provide a reliable and unambiguous test of the predictions of ChPT.

In our previous papers \cite{15–17} we have formulated a consistent field-theoretical approach to the calculation of the pionium observables. Below the basic ideas and assumptions of our approach are given.

The pions within our approach are described by the elementary fields in the Lagrangian. This allows one to write down the bound-state BS equation for the metastable $\pi^+\pi^-$ atom. The necessary link to the strong amplitudes from ChPT where the pions are described by the pseudoscalar quark densities, proceeds via the Deser-type factorization in analogy with the potential theory. Namely, we express all corrections to the Deser formula in terms of the on-mass-shell pion amplitudes which are further identified with the amplitudes from ChPT.

In our calculations we always adopt the so-called ”local” approximation which consists is suppressing the relative momentum dependence of the strong amplitudes. The origin of this approximation can be traced back to the huge difference between the bound-state and strong momentum scales. Since the tree-level strong amplitude for the process $\pi^+\pi^- \to \pi^0\pi^0$ does not depend on the relative momentum of the pions, the corrections arising due to this approximation, by a simple power counting, should be suppressed by a product of two small factors: the fine structure constant $\alpha$ which is the ratio of the bound-state momentum and the pion mass, and the factor $m_\pi/(4\pi F_\pi) \sim 10^{-1}$ being the ratio of two strong scales in chiral theories. For this reason we at the present stage completely neglect the corrections coming from using the ”local” approximation.

The pionium in our approach is described by the exact bound-state BS equation

$$\langle \chi | G^{-1}_0 (P) = \langle \chi | V (P) \right.$$  \hspace{1cm} (2)

where $\chi$ denotes the BS wave function of pionium, $G_0$ is the free Green’s function of the $\pi^+\pi^-$ pair and $V$ is the full BS kernel which apart from all two-particle irreducible diagrams includes the self-energy corrections in two outgoing charged pion legs \cite{24,25}. The square of the c.m. momentum $P$ takes the value $P^2 = M^2 = M^2 - iM\Gamma$ where $M$ denotes the ”mass” of an atom, and $\Gamma$ stands for the decay width.

In order to perform the perturbative expansion of the bound-state observables we single out from the full kernel $V$ the instantaneous Coulomb part $V_C$ which is responsible for the formation of the bound state composed of $\pi^+$ and $\pi^-$. The ”remainder” of the potential which is treated perturbatively is denoted by $V' = V - V_C$.

The ”unperturbed” part of the full kernel $V_C$ is chosen in the form guaranteeing the exact solvability of the corresponding BS equation. The solution of this equation yields the
relativistic Coulomb wave function $\psi_C$ with the corresponding eigenvalue $P^{*2} = M^{*2} = m_{\pi}^2(4 - \alpha^2)$ where $m_{\pi}$ is the charged pion mass. The solution of the full BS equation is expressed in terms of the unperturbed solution as follows [16, 17]

$$<\chi| = \text{const}\times<\psi_C|[1 + (G_0^{-1}(P) - G_0^{-1}(P^*) - V'(P))G_RQ]^{-1}$$

where $G_R$ denotes the pole-subtracted part of the exact relativistic Coulomb Green’s function and $Q$ is the projection operator on the subspace orthogonal to the unperturbed ground-state BS wave function [16, 17]. Substituting this solution into the complete BS equation (2) we arrive at the final relation

$$<\psi_C|[1 + (G_0^{-1}(P) - G_0^{-1}(P^*) - V'(P))G_RQ]^{-1}(G_0^{-1}(P) - G_0^{-1}(P^*) - V'(P))]|\psi_C > = 0$$

which serves as the basic equation for performing the perturbative expansion of the bound-state observables. The explicit form of the operator $G_RQ$ is given and the perturbation kernel $V'$ is known up to any given order of loop expansion. The only unknown quantities entering (parametrically) into this expression are the real and imaginary parts of $P^2$. Consequently, expanding this expression in powers of $V'$, we arrive at the recursive relations which define the mass and width of the bound state in any given order.

In order to provide the necessary Deser-type factorization in the obtained results, one needs a further classification of diagrams entering into the definition of $V'$. We consider the following subsets of diagrams:
1. A purely strong part, which is isotopically invariant. This part survives when electromagnetic interactions are ”turned off” the theory.
2. The part which is responsible for the $m_{\pi^0} - m_{\pi^0}$ electromagnetic mass difference
3. Remaining electromagnetic effects, including the exchanges of virtual photons.

The part 1 includes strong interactions which govern the decay of a pionium. The part 2 makes this decay kinematically allowed. Consequently, it is natural to consider them together, denoting the corresponding potential as $V_{12} = V_1 + V_2$. The $T$-matrix corresponding to summation of the potential $V_{12}$ in all orders is given by $T_{12}(P) = V_{12}(P) + V_{12}(P)G_0(P)T_{12}(P)$. The rest of the potential is referred to as $V_3 = V' - V_{12}$ and is treated perturbatively. Further, the quantity $G_RQ$ in our basic relation can be written as $G_RQ = G_0(P^*) + \delta G$ where $\delta G$ corresponds to the exchange of the ladder of Coulomb photons and is also considered as a perturbation. Then in the first order in $V_3$ and $\delta G$ the basic relation takes the form

$$0 = -2iM^*\delta M - <\psi_C|T_{12}|\psi_C > +$$

$$+ \delta M <\psi_C|(G_0^{-1})'G_0T_{12}|\psi_C > + (\delta M)^2 <\psi_C|(1 + T_{12}G_0)(G_0^{-1})''|\psi_C > +$$

$$+ <\psi_C|\delta M(G_0^{-1})' - T_{12})\delta GT_{12}|\psi_C > - <\psi_C|(1 + T_{12}G_0)\delta G|\psi_C >$$

where $G_0 = G_0(M^*)$ and the prime stands for the differentiation with respect to $M^*$.

In the ”local” approximation $T_{12}$ does not depend on the relative momenta and thus can be taken out of the matrix elements. Neglecting first the corrections (only the first line of Eq. (3) contributes), we arrive at the relativistic counterpart of Eq. (4)

$$\Delta E^{(1)} = \text{Re} \left( \frac{i T_{12}}{2M^* m_{\pi} \phi_0^2} \right), \quad -\frac{1}{2} \Gamma^{(1)} = \text{Im} \left( \frac{i T_{12}}{2M^* m_{\pi} \phi_0^2} \right)$$

(5)
\[
\text{Re} \, (iT_{12}) = 16\pi \, T_{\pi^+\pi^- \to \pi^+\pi^-} (4m_{\pi}^2; \vec{0}, \vec{0}) \]
\[
\text{Im} \, (iT_{12}) = -16\pi \left( \frac{\Delta m_{\pi}}{2m_{\pi}} \right)^{1/2} \left( 1 - \frac{\Delta m_{\pi}}{2m_{\pi}} \right)^{1/2} |T_{\pi^+\pi^- \to \pi^0\pi^0}(4m_{\pi}^2, \vec{0}, \vec{q}_0)|^2,
\]

and \( T(s; \vec{p}, \vec{q}) \) denote the (dimensionless) \( S \)-wave \( \pi \pi \) strong scattering amplitudes which include the effect of \( m_{\pi^\pm} - m_{\pi^0} \) mass difference. \( \vec{q}_0 \) is the relative momentum of the \( \pi^0\pi^0 \) pair at the threshold \( s = 4m_{\pi}^2 \), with the magnitude given by the relation \( m_{\pi}^2 = m_{\pi^0}^2 + \vec{q}_0^2 \).

In order to calculate the corrections to the relativistic Deser formula, one has to evaluate the integrals entering into the Eq.\( (\ref{eq:correction}) \).

1. The correction due to the shift of the bound-state pole by strong interactions is determined by two terms in the second line of Eq.\( (\ref{eq:correction}) \). Calculating these integrals explicitly and using Eqs.\( (\ref{eq:correction}) \) we obtain

\[
\delta_S = -\frac{9}{8} \frac{\Delta E^{(1)}}{E_1} = -5.47 \times 10^{-3} \, m_{\pi}(2a_0^0 + a_0^2)
\]

where \( E_1 \) denotes the unperturbed ground-state binding energy.

2. The correction due to the relativistic modification of the Coulomb wave function is given by the value of the wave function \( \psi_C \) at the origin which for the particular choice of the kernel used in the calculations is related to its nonrelativistic counterpart by

\[
\psi_C(0) = \Psi_1(0) \left( 1 - 0.381\alpha + \cdots \right)
\]

and the correction in the decay width which is proportional to \( |\psi_C(0)|^2 \) is twice as large.

3. The correction due to the exchange of Coulomb photon ladders is given by the first term in the third line of Eq.\( (\ref{eq:correction}) \) containing \( \delta G \). Using the known expression for the relativistic Coulomb Green’s function\( \ref{eq:G_Coulomb} \)\( \ref{eq:G_Coulomb} \) and evaluating this integral explicitly, we obtain

\[
\delta_C = (1/2 + 2.694 - \ln\alpha) \, \Delta E^{(1)}/E_1 = 3.95 \times 10^{-2} \, m_{\pi}(2a_0^0 + a_0^2)
\]

4. In order to obtain the mass shift and radiative corrections, in the perturbation kernel \( V_3 \) one has to take into account the residual photon exchange diagram and the self-energy corrections in the outgoing charged pion legs. The net result of this effect consists in the cancellation of the correction due to the relativistic modification of the Coulomb wave function and in replacing the quantity \( T_{\pi^+\pi^- \to \pi^0\pi^0} \) in Eq.\( (\ref{eq:correction}) \) by the Coulomb pole removed full amplitude for the process \( \pi^+\pi^- \to \pi^0\pi^0 \). Using the expression of this amplitude\( \ref{eq:G_Coulomb} \) and isolating the purely strong isotopically symmetric amplitude with a common mass equal to the charged pion mass, we arrive at the following expressions for the mass shift and radiative corrections

\[
\delta_M = \frac{2\Delta m_{\pi}^2}{3m_{\pi}^2} \left( 1 + \frac{m_{\pi}^2}{96\pi^2 F_{\pi}^2} \left( 48 + \frac{16}{3} l_1 - \frac{16}{3} l_2 + 7l_3 - 36l_4 \right) \right)
\]

\[
\delta_{\text{em}} = \frac{\alpha}{12\pi} (-30 + 3\mathcal{K}_{1}^{\pm 0} - \mathcal{K}_{2}^{\pm 0})
\]

where \( \mathcal{K}_{i}^{\pm 0} \) and \( l_i \) denote the low-energy constants of ChPT\( \ref{eq:K_i} \)\( \ref{eq:K_i} \).

5. The correction due to the vacuum polarization is of order \( O(\alpha^2m_{\pi}/m_{e}) \) (the contribution due to the hadronic vacuum polarization is negligibly small). Taking into account the
corresponding diagram in the kernel $V_3$ and evaluating numerically the resulting integral, we arrive at the following result

$$\delta_{\text{vac}} = \frac{3}{16} \alpha^2 \frac{m_\pi}{m_e} \times 0.6865$$  \hspace{1cm} (13)$$

which completely agrees with the zero-Coulomb piece of the result given in Ref. 18.

The additional contributions coming from one-Coulomb and many-Coulomb pieces which are also considered in this paper, within our approach arise in the second-order of the perturbation theory.

6. The correction to the pionium decay width due to the finite size effect is caused by the modification of the instantaneous Coulomb interaction by the pion loop in the two pion-photon vertex (the corresponding diagram is contained in the kernel $V_3$). For simplicity, we approximate this vertex by a monopole parameterization. After calculating the corresponding matrix element explicitly, we obtain

$$\delta_F = \frac{2\alpha}{3\pi} m_\pi^2 < r^2 >_V \ln \left( \frac{1}{24} m_\pi^2 < r^2 >_V \right)$$ \hspace{1cm} (14)$$

where $< r^2 >_V$ denotes the square charge radius of the pion.

The relativistic Deser formula with all above corrections takes the form

$$\tau_1^{-1} = \frac{16\pi}{9} \left( \frac{2 \Delta m_\pi}{m_\pi} \right)^{1/2} \left( 1 - \frac{\Delta m_\pi}{2m_\pi} \right)^{1/2} (a_0^0 - a_0^2)^2 \phi_0^2 (1 + \delta_S + \delta_C + \delta_M + \delta_{\text{em}} + \delta_{\text{vac}} + \delta_F)$$

where $a_0^0$ and $a_0^2$ denote the $\pi\pi$ scattering lengths in the isospin-symmetric case, with the charged pion mass taken to be the common mass of the pion octet.

In our calculations for the constants $\tilde{l}_i$ we take the numerical values from Ref. 17

$$\tilde{l}_1 = -2.3 \pm 2.7, \quad \tilde{l}_2 = 6.0 \pm 1.3, \quad \tilde{l}_3 = 2.9 \pm 2.4, \quad \tilde{l}_4 = 4.3 \pm 0.9.$$  

Also, we use the values $\frac{e^2 F^2}{m_\pi^2} K_{1}^{\pm 0} = 1.8 \pm 0.9, \quad \frac{e^2 F^2}{m_\pi^2} K_{2}^{\pm 0} = 0.5 \pm 2.2$ from Ref. 20. Other input parameters in our calculations are the S-wave $\pi\pi$ scattering lengths: $a_0^0 = 0.217 m_\pi^{-1}, a_0^2 = -0.041 m_\pi^{-1}$ calculated in ChPT and the e.m. charge radius of pion $< r^2 >_V = 0.439 Fm$.

The results of our calculations are shown in Table 1.

As we observe from Table 1, the largest correction in the decay width is caused by the mass shift effect. Our result for the sum of mass shift and electromagnetic radiative corrections generally agrees in sign and magnitude with the recent field-theoretical calculations from Ref. 14.

The sign of the mass shift effect obtained in the nonrelativistic scattering theory approach 19 turns out to be opposite as compared to our result, and is of the same order of magnitude.

Our result for the correction due to the exchange of Coulomb photons agrees with the result obtained in the potential scattering theory 19 and disagrees with the result from Refs. 13, 14. Numerically the largest part in this effect comes from the nonanalytic ln piece which is exactly the same in our approach and in the potential theory, and is absent in Refs. 13, 14.

Thus, we have evaluated a complete set of the lowest-order corrections to the pionium decay width. For a full understanding of the problem, however, the reason for the difference in sign in the mass shift effect in the field-theoretical and potential theories should be investigated in detail. Also, one should have a reliable quantitative estimate of the accuracy of ”local” approximation including a proper treatment of arising new UV divergences.
Table 1. Corrections to the $\pi^+\pi^-$ atom decay width.

| Effect                         | Value | Correction (in %) |
|-------------------------------|-------|-------------------|
| Strong                        | $\delta_S$ | -0.22             |
| Coulomb photon exchange       | $\delta_C$ | +1.55             |
| Mass shift                    | $\delta_M$ | +2.99 ± 0.77      |
| Electromagnetic radiative     | $\delta_{em}$ | +1.73 ± 2.31     |
| Vacuum polarization           | $\delta_{vac}$ | +0.19          |
| Finite size                   | $\delta_F$ | -0.16             |
| Total                         | $\delta_{tot}$ | +6.1 ± 3.1      |

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References

1. J. Gasser and H. Leutwyler, *Ann. Phys. (N.Y.*)** 158, 142 (1984).
2. J. Bijens et al., *Phys. Lett. B** 374, 210 (1996).
3. M. Knecht et al., *Nucl. Phys. B** 457, 513 (1995); *B** 471, 445 (1996).
4. D. Počanić, Talk given at the Workshop on Chiral Dynamics: Theory and Experiment, Mainz, Germany, 1-5 Sep 1997, [hep-ph/9801360]; D. Počanić, these proceedings.
5. B. Adeva et al., proposal to the SPSLC, CERN/SPSLC 95-1 (1995).
6. S. Deser et al., *Phys. Rev.** 96, 774 (1954).
7. G.V. Efimov, M.A. Ivanov and V.E. Lyubovitskij, *YAF** 44, 460 (1986); *Pis’ma v JETP** 45, 526 (1987).
8. A.A. Bel’kov, V.N. Pervushin and F.G. Tkebuchava, *YAF** 44, 466 (1986).
9. A. Karimkhodzaev and R.N. Faustov, Dubna preprint P2-86-142.
10. M.K. Volkov, *Teor. Mat. Fiz.*, 71, 381 (1987).
11. Z. Silagadze, *JETP Lett.*, 60, 689 (1994).
12. E.A. Kuraev, *Part. Nucl. Phys.*, 61, 239 (1998).
13. H. Jallouli and H. Sazdjian, *Phys. Rev.** D58, 014011 (1998).
14. H. Sazdjian, these proceedings.
15. V. E. Lyubovitskij and A. G. Rusetsky, *Phys. Lett. B** 389, 181 (1996);
16. V. E. Lyubovitskij, E. Z. Lipartia and A. G. Rusetsky, *JETP Lett.*, 66, 747 (1997);
17. M. A. Ivanov et al., *hep-ph/9805356*, *Phys. Rev.** D58, in print.
18. P. Labelle and K. Buckley, Preprint *hep-ph/9804201*.
19. U. Moor, G. Rasche and W.S. Woolcock, *Nucl. Phys.** A587, 747 (1995); A. Gashi, G. Rasche, G. C. Oades and W. S. Woolcock, *Nucl. Phys.** A628, 101 (1998); A. Gashi and G. Rasche, these proceedings.
20. M. Knecht and R. Urech, *Nucl. Phys. B** 519, 329 (1998).