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PIONIC EXCITATIONS AND CHIRAL SYMMETRY IN DENSE MATTER* **

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This presentation discusses some recent developments in our understanding of the lightest quark–antiquark excitations, both of the condensed QCD vacuum and in a nuclear medium. We focus on two selected topics: the thermodynamics of the chiral condensate and in-medium s-wave interactions of pions with special emphasis on the recently observed deeply bound pionic atom states.

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1. Prelude: symmetries and symmetry breaking patterns in QCD

The QCD ground state, or vacuum, is characterized by the presence of a strong condensate $\langle \bar{q}q \rangle$ of scalar quark–antiquark pairs (the chiral condensate) which represents the order parameter for spontaneous chiral symmetry breaking in QCD. The light hadrons are quasi-particle excitations of this condensed ground state. Pions and kaons are of special importance in this context, as they are identified with the pseudoscalar Goldstone bosons of spontaneously broken chiral symmetry; their masses would vanish in the limit of massless u-, d- and s-quarks. The pion decay constant, $f_\pi \simeq 92.4$ MeV, determines the chiral scale $4\pi f_\pi \sim 1$ GeV (refered to as the "chiral gap") which governs the low-mass hadron spectrum. For example, the lightest vector mesons (ρ, ω) can be interpreted as the lowest resonant $q\bar{q}$ "dipole" excitations of the QCD vacuum. Current algebra combined with QCD finite energy sum rules [1, 2] connects their masses directly with the chiral gap,

$$\sqrt{2}\, m_V = 4\pi f_\pi \,, \tag{1}$$

in leading order.

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The deviation of the physical pion mass, $m_\pi \simeq 0.14$ GeV, from zero reflects weak explicit chiral symmetry breaking by the small masses of u- and d-quarks, $m_{u,d} < 10$ MeV. The larger kaon mass ($m_K \simeq 0.5$ GeV) shows the stronger explicit symmetry breaking caused by the mass of the strange quark, $m_s \sim 0.15$ GeV. Spontaneous and explicit chiral symmetry breaking imply the PCAC or Gell–Mann, Oakes, Renner (GOR) relation,

$$m_\pi^2 f_\pi^2 = -\tfrac{1}{2}(m_u + m_d)\langle \bar{u}u + \bar{d}d \rangle \ , \tag{2}$$

to leading order in the quark masses $m_{u,d}$.

One of the basic issues in strong interaction physics is to explore the QCD phase diagram as it evolves with increasing temperature and/or baryon chemical potential. A key element in this discussion is the chiral transition from the Nambu–Goldstone realization of chiral symmetry (with non-zero condensate $\langle \bar{q}q \rangle$) to the "restored" Wigner–Weyl realization in which the chiral condensate vanishes. In QCD, chiral restoration is probably linked to the transition between composite hadrons and deconfined quarks and gluons.

Chiral perturbation theory [3] determines the (weak) leading temperature dependence of the condensate for N_f massless quark flavours as

$$\frac{\langle \bar{q}q \rangle_T}{\langle \bar{q}q \rangle_0} = 1 - \frac{N_f^2 - 1}{N_f}\left(\frac{T^2}{12 f_\pi^2}\right)\left(1 + \frac{T^2}{24 N_f f_\pi^2}\right) + O(T^6) \ . \tag{3}$$

Lattice QCD [4] locates the critical temperature for the chiral transition at $T_c \simeq$(150–200) MeV. The leading dependence of $\langle \bar{q}q \rangle$ on baryon density ρ at zero temperature is controled by the pion–nucleon sigma term, $\sigma_N \simeq$ 0.5 GeV:

$$\frac{\langle \bar{q}q \rangle_\rho}{\langle \bar{q}q \rangle_0} \simeq 1 - \frac{\sigma_N}{m_\pi^2 f_\pi^2}\rho + \dots \ , \tag{4}$$

indicating a rapidly decreasing magnitude of the chiral condensate in cold compressed nuclear matter [5].

The GOR relation (1) continues to hold [6] in matter at finite temperature $T < T_c$ and at finite density, when reduced to a statement about the *time* component, A_0, of the axial current A_μ. We denote the in-medium pion decay constant related the thermal matrix element $\langle \pi | A_0 | \rangle_{\rho,T}$ by $f_\pi^*(\rho, T)$. One finds

$$f_\pi^{*2}(\rho, T) = -\frac{m_q}{m_\pi^{*2}}\langle \bar{q}q \rangle_{\rho,T} + \dots, \tag{5}$$

to leading order in the average quark mass $m_q = 1/2\ (m_u + m_d)$, where $\langle \bar{q}q \rangle_{\rho,T}$ stands for the T- and ρ-dependent condensate, $\langle \bar{u}u + \bar{d}d \rangle_{\rho,T}$. The "melting" of the condensate by heat or compression therefore translates primarily into an in-medium change of the decay constant of the pion, given that its mass, m_π^*, is not much affected by the medium because of its Goldstone boson nature.

The fact that the "chiral gap", $4\pi f_\pi^*(\rho,T)$, decreases when thermodynamic conditions change toward chiral restoration, should also imply characteristic observable changes in the low-energy, s-wave dynamics of pions in matter, and in the meson mass spectrum.

The emphasis in this presentation will be on two selected topics of current interest. First we study the thermodynamics of the chiral condensate, $\langle\bar{q}q\rangle_{\rho,T}$. Then we discuss cold matter under moderate conditions and examine the influence of a density-dependent pion decay constant, $f_\pi^*(\rho)$, on the recently observed deeply bound pionic atomic states in Pb.

2. Thermodynamics of the chiral condensate

2.1. Basics

Suppose we are given a chiral effective Lagrangian, $\mathcal{L}_{\rm eff}$, with Goldstone bosons (pions) coupled to baryons (nucleons). Let $\mathcal{Z}$ be the partition function derived from this theory, and μ the baryon chemical potential. The pressure as a function of μ and T is

$$P(\mu,T) = \frac{T}{V}\ln\mathcal{Z}\ , \tag{6}$$

where V is the volume. Note that the Hamiltonian which determines $\mathcal{Z}$ depends on the pion mass m_π, or equivalently, on the quark mass m_q through the GOR relation (2). Given the equation of state $P(\mu,T)$, a variant of the Hellmann–Feynman theorem (with the quark mass treated formally as an adiabatic parameter) leads to the following expression for the density and temperature dependent chiral condensate:

$$\langle\bar{q}q\rangle_{\rho,T} = \langle\bar{q}q\rangle_0 - \frac{dP(\mu,T)}{dm_q}\ , \tag{7}$$

where the baryon density is $\rho = \partial P/\partial\mu$. Using the GOR relation (1), one can rewrite (7) as:

$$\frac{\langle\bar{q}q\rangle_{\rho,T}}{\langle\bar{q}q\rangle_0} = 1 + \frac{1}{f_\pi^2}\frac{dP(\mu,T)}{dm_\pi^2}\ . \tag{8}$$

(For further details see *e. g.* [7].) The task is therefore to investigate how the equation of state, at given temperature and baryon chemical potential, changes when varying the squared pion mass (or the quark mass).

Note that for a system of nucleons (mass M) interacting with pions, the total derivative of the pressure with respect to m_π^2 reduces to

$$\frac{dP(\mu,T)}{dm_\pi^2} = \frac{\partial P(\mu,T)}{\partial m_\pi^2} - \frac{\sigma_N}{m_\pi^2}\rho_S(\mu,T) \ , \tag{9}$$

with the scalar density $\rho_S = -\partial P/\partial M$ and the sigma term $\sigma_N = m_q \partial M/\partial m_q$.

2.2. Model

Consider now the following effective Lagrangian as an approximation to the hadronic phase of QCD:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_N + \mathcal{L}_\pi + \mathcal{L}_{\pi N} + \mathcal{L}_{NN} \ . \tag{10}$$

The free nucleon Lagrangian is $\mathcal{L}_N = \bar{N}(i\gamma \cdot p - M)N$, where M is the nucleon mass in vacuum. The pion sector with inclusion of $\pi\pi$ interactions is described by the non-linear sigma model plus the standard pion mass term:

$$\mathcal{L}_\pi = \frac{f_\pi^2}{4}\,\mathrm{tr}\,[\partial_\mu U \partial^\mu U^+] + \text{ mass term} \,, \tag{11}$$

with $U = \exp[i\vec{\tau}\cdot\vec{\pi}/f_\pi]$. The chiral pion–nucleon coupling to leading order in pion momentum is

$$\mathcal{L}_{\pi N} = \frac{g_A}{2f_\pi}\bar{N}\gamma_\mu\gamma_5\vec{\tau}N \cdot \partial^\mu\vec{\pi} - \frac{1}{4f_\pi^2}\bar{N}\gamma_\mu\vec{\tau}N \cdot \vec{\pi} \times \partial^\mu\vec{\pi} \ . \tag{12}$$

The short-distance dynamics is absorbed in NN contact terms,

$$\mathcal{L}_{NN} = -\frac{G_S}{2}(\bar{N}N)^2 + \frac{G_V}{2}(\bar{N}\gamma_\mu N)^2 + \ldots \,, \tag{13}$$

with the coupling strength parameters $G_{S,V}$ fixed by the ground state properties of normal nuclear matter. What we have in mind in the first step is a variant of relativistic mean field theory combined with "soft" pion fluctuations treated within the framework of chiral perturbation theory.

We have used two-loop thermal field theory to perform a self-consistent calculation of $P(\mu,T)$ and then deduced the chiral condensate as a function of temperature and baryon density using Eq. (8). This calculation [8] generates temperature dependent mean fields for the nucleons at the same time as it treats thermal pion fluctuations with inclusion of leading $\pi\pi$ interactions. The pressure equation takes the form:

$$P(\mu,T) = P_N(\mu^*, M^*, T) + P_\pi(\mu^*, M^*, T) + \frac{G_V}{2}\rho^2 - \frac{G_S}{2}\rho_S^2 \ , \tag{14}$$

with nucleon and pion contributions $P_{N,\pi}$, respectively, depending on the effective nucleon mass

$$M^* = M - G_S \rho_S \tag{15}$$

and the shifted baryon chemical potential

$$\mu^* = \mu - G_V \rho \; . \tag{16}$$

The baryon and scalar densities are determined as

$$\rho = \frac{\partial}{\partial \mu^*}(P_N + P_\pi), \; \rho_S = \frac{\partial}{\partial M^*}(P_N + P_\pi) \; . \tag{17}$$

The whole set of equations (14)–(17) is then solved self-consistently.

2.3. Quark condensate at finite T and ρ

We return now to Eq. (8) and evaluate the variation of the chiral condensate with temperature and density. The dependence of $P(\mu, T)$ on the pion mass is explicit in the thermal pion Green function and implicit through the nucleon mass, using Eq. (9). The result for $\langle \bar{q}q \rangle_{\rho,T}$ is shown in Fig 1. One notes that the temperature dependence at $\rho = 0$ is quite similar to the result of lattice QCD [4] and close to the one found in chiral perturbation theory [3]. The critical temperature in the present calculation is $T_c \simeq 180$ MeV.

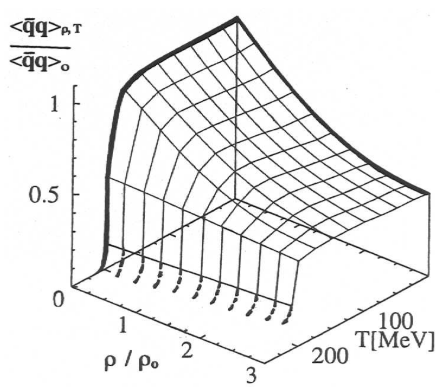

Fig. 1. Dependence of the chiral condensates $\langle \bar{q}q \rangle_{\rho,T}$ on temperature T and baryon density ρ, calculated [8] in two-loop thermal field theory with the effective Lagrangian (10)–(13).

At $T = 0$ we have $m_q(\langle \bar{q}q \rangle_\rho - \langle \bar{q}q \rangle_0) = \sigma_N \rho_S$. At low baryon density where $\rho_S \simeq \rho$ the linear behaviour as in Eq. (4) is recovered. Naive extrapolation of this linear density dependence with $\sigma_N \simeq 45$ MeV would find the condensate dropping to zero at about three times nuclear matter density $\rho_0 \simeq 0.17$ fm^{-3}. However the scalar density ρ_S becomes significantly smaller than ρ at high density, so that $\langle\bar{q}q\rangle_\rho$ still keeps almost half of $\langle\bar{q}q\rangle_0$ at $\rho = 3\rho_0$. For densities up to ρ_0, however, the linear behaviour (4) turns out to be a good approximation. Pionic fluctuations, though not of great overall importance at $T = 0$, help to maintain this simple linear dependence.

Our calculations are presently extended to the three-loop level, so we will see whether the picture, Fig 1, is going to persist (see also Ref. [9]). Up to this point we can conclude that the magnitude of the quark condensate at $\rho \simeq \rho_0$ is reduced by about one third from its vacuum value, so that this tendency toward "chiral restoration" should have observable consequences already in normal nuclear systems.

3. Pionic s-waves in the nuclear medium

The investigation of pion–nucleus interactions has a long history [10,11]. The reasons for revisiting this topic are two-fold: first, the recent observation of deeply bound pionic atom states in Pb isotopes [12] has sharpened the quantitative constraints on the local (s-wave) part of the pion–nuclear optical potential, and secondly, there is renewed interest in the theoretical foundations of this optical potential from the point of view of chiral dynamics [13]. In the present context we concentrate primarily on this latter point.

Pions as Goldstone bosons interact weakly at low momentum. Their s-wave interactions with nucleons in leading order are determined by the pion decay constant f_π as the relevant scale of spontaneously broken chiral symmetry. In the nuclear medium, this scale changes, and the obvious question is whether accurate data, such as those from deeply bound pionic atoms, are a sensitive measure for the expected density dependence of f_π.

3.1. Chiral pion–nuclear dynamics

The spectrum of pionic modes with energy ω and momentum $\vec{q}$ in nuclear matter at density ρ is determined by solutions of the wave equation

$$\left[\omega^2 - \vec{q}^{\,2} - m_\pi^2 - \Pi(\omega, \vec{q}; \rho)\right] \phi = 0 \,, \tag{18}$$

where the self-energy Π is often expressed in terms of the optical potential U as $\Pi = 2\omega U$. For pions "at rest" relative to the surrounding matter, it is convenient to introduce an in-medium effective mass [13] by $m_\pi^{*2}(\rho) \equiv \omega^2(\vec{q} = 0; \rho) = m_\pi^2 + \mathrm{Re}\,\Pi(\omega = m_\pi^*, \vec{q} = 0; \rho)$. The mass shift in matter is then a measure of the underlying s-wave pion–nuclear interactions.

Consider now a low-energy π^- interacting with matter at low proton and neutron densities $\rho_{p,n}$. To leading order in these densities,

$$\Pi = 2\omega U = -T(\pi^- p)\rho_p - T(\pi^- n)\rho_n \ , \tag{19}$$

where T denotes the πN T-matrix (at threshold, its relation to the corresponding scattering length is $T(\vec{q}=0) = 4\pi(1 + m_\pi/M)a$).

The low-energy behaviour of T is ruled by theorems based on chiral symmetry. Consider the isospin even and odd amplitudes, $T^{(\pm)} = 1/2\,[T(\pi^- p) \pm T(\pi^- n)]$. The Tomozawa–Weinberg theorem gives

$$T^{(+)}(\omega, \vec{q}=0) = 0, \ \ T^{(-)}(\omega, \vec{q}=0) = \frac{\omega}{2f_\pi^2} \ , \tag{20}$$

to leading order in ω. In next-to-leading order an attractive scalar term proportional to σ_N/f_π^2 combines with a repulsive term of order ω^2 so as to reproduce the observed very small isospin-averaged scattering length, $a^{(+)} = (-0.003 \pm 0.002)$ fm. For the case of $T^{(-)}$, chiral perturbation theory gives corrections of order ω^3 which close the 15 % gap between the lowest order result (20) and the empirical isospin-odd scattering length, $a^{(-)} = (0.128 \pm 0.002)$ fm. (See Ref. [14] for a recent analysis of πN scattering lengths.) In the actual calculations we use the threshold amplitudes $T^{(+)} = 0$ and $T^{(-)} = m_\pi/2f_\pi^2(1 + 0.066\, m_\pi^2/f_\pi^2)$, compatible with the empirical scattering lengths.

To leading chiral order and in the low-density limit, the π^- self-energy is simply $2\omega U(\pi^-) = T^{(-)}(\rho_n - \rho_p)$, with $T^{(-)}$ given by (20). For a π^+ the isospin-odd part of the amplitude changes sign, so that the primary medium effect is a splitting of the π^+ and π^- masses in asymmetric nuclear matter:

$$\Delta m(\pi^\pm) = U(\pi^\pm) = \pm \frac{\rho_p - \rho_n}{4f_\pi^2} \, . \tag{21}$$

In isospin-symmetric matter there is no shift of the pion mass to this order. Clearly, systems with a large neutron excess are of special interest here.

The s-wave pion–nucleus optical potential involves more than just the amplitudes in leading order. Double scattering terms are known to be important, and absorptive corrections of order ρ^2 must be added [10, 11]. The s-wave in-medium self-energy for a π^- becomes

$$\Pi(\vec{q}=0) = -T_{\text{eff}}^{(+)}(\rho_p + \rho_n) - T_{\text{eff}}^{(-)}(\rho_p - \rho_n) - 4\pi B_0 \rho^2 \ , \tag{22}$$

where T_{eff} includes double scattering effects. With $T^{(+)} = 0$ one finds [10,11] in the Fermi gas approximation:

$$T_{\text{eff}}^{(+)} = -\frac{3p_{\text{F}}}{4\pi^2} T^{(-)^2} \ , \qquad T_{\text{eff}}^{(-)} = T^{(-)}\left(1 - \frac{3p_{\text{F}}}{8\pi^2} T^{(-)}\right) , \tag{23}$$

where $p_{\rm F}$ is the nuclear Fermi momentum. Note that, with $T^{(+)} = 0$, the leading term in $T^{(+)}_{\rm eff}$ now involves the squared isospin-odd amplitude proportional to f_π^{-4}. In $T^{(-)}_{\rm eff}$ the double scattering correction (about -10% at $p_{\rm F} \simeq 2m_\pi$) is often ignored, but we prefer to take it into account[1]. The ρ^2-term has a complex, phenomenological constant B_0. Its imaginary part is fitted to reproduce pion absorption rates.

Apart from the $B_0\rho^2$ term, the s-wave π^- potential at $\omega = m_\pi$ and for a given N/Z ratio $\eta = \rho_n/\rho_p$ becomes

$$U(\vec{q}=0) \simeq 8 \text{ MeV}\left(\frac{\rho}{\rho_0}\right)^{4/3} + 44 \text{ MeV}\left(\frac{\eta-1}{\eta+1}\right)\frac{\rho}{\rho_0}\left[1-0.1\left(\frac{\rho}{\rho_0}\right)^{1/3}\right], \quad (24)$$

which gives about 16 MeV of repulsion at $\rho = \rho_0$ and $\eta = 1.5$, the N/Z ratio characteristic of the Pb region. From the analysis of pionic atom data it is known that this repulsion is too weak by about a factor of two. It has been common practice to choose the phenomenological ReB_0 such that the missing repulsion is accounted for. This requires a large negative ReB_0 for which there is little theoretical foundation. In fact, the two-body mechanisms which are held responsible for ReB_0 should also be present in the deuteron. But the real part of the $\pi^- d$ scattering length is perfectly well reproduced just by the single and double scattering terms already present in $T_{\rm eff}$, suggesting Re$B_0 \simeq 0$. We should look for an alternative way to generate the "missing repulsion".

3.2. In-medium chiral condensate and ρ dependent pion decay constant

The considerations in the previous section were based on the chiral low-energy theorem (20), expressed in terms of the *vacuum* pion decay constant. In other words: our reference point for the chiral expansion (in powers of the pion energy or momentum) has so far been the minimum of the vacuum effective potential derived from the chiral Lagrangian, the one that applies at density $\rho = 0$. However, the nuclear medium defines a *new* vacuum, with the minimum of the effective potential shifted such that the magnitude of the chiral condensate $\langle\bar{q}q\rangle$ is reduced (see Section 2). Following the in-medium GOR relation (5) for the *time* component of the axial current, we have $\langle\bar{q}q\rangle_\rho/\langle\bar{q}q\rangle_0 = f_\pi^{*2}(\rho)/f_\pi^2$ (assuming at this stage no change of the Goldstone boson mass). The shift of the vacuum therefore implies a density dependent pion decay constant,

$$f_\pi^{*2}(\rho) = f_\pi^2 - \frac{\sigma_N}{m_\pi^2}\rho\,, \quad (25)$$

to leading order in the baryon density.

[1] In the literature the notations $T^{(+)}_{\rm eff} = 4\pi(1+m_\pi/M)b_0^{\rm eff}$ and $T^{(-)}_{\rm eff} \simeq T^{(-)} = -4\pi(1+m_\pi/M)b_1$ are frequently used.

While the minimum of the effective potential is shifted at $\rho > 0$, the chiral low-energy theorem (20) for πN scattering still holds, but now with respect to the *new* vacuum with its reduced condensate and reduced pion decay constant $f_\pi^*(\rho)$. The isospin-odd in-medium πN amplitude becomes

$$T^{(-)} = \frac{\omega}{2f_\pi^{*2}}\,, \tag{26}$$

to leading chiral order, while the isospin-even amplitude still has $T^{(+)} = 0$. Chiral perturbation theory should now be expanded around the new, shifted vacuum, with the chiral gap replaced by $4\pi f_\pi^*(\rho)$ wherever it appears.

With Eq. (25), this immediately implies that the s-wave optical potential (22), with f_π replaced by f_π^*, will be substantially more repulsive than the one given by Eq. (24). In fact, this potential becomes about twice as large at $\rho = \rho_0$ and $\eta = 1.5$ when replacing f_π by $f_\pi^*(\rho_0) \simeq 0.82 f_\pi$ according to Eq. (25). We will now investigate the consequences of this assertion for the understanding of deeply bound pionic atom states in heavy nuclei.

3.3. Deeply bound pionic states

The existence of narrow $1s$ and $2p$ states in heavy pionic atoms results from a subtle balance between the attractive Coulomb potential and the repulsive s-wave optical potential [15]. The net attraction is localized at and beyond the nuclear surface. Under these special conditions the overlap of the pion densities with the nuclear density distribution is sufficiently small so that the absorptive width is reduced and the deeply bound states have a chance to be observed as narrow structures. This is the case in the GSI measurements of $1s$ and $2p$ pionic states in ^{207}Pb [12] and, most recently, in ^{205}Pb, using the (d, ^{3}He) reaction for their production.

We have performed detailed calculations for pionic $1s$ and $2p$ states in Pb and other isotopes [16]. The aim is to explore, in particular, the sensitivity of the widths of these states with respect to the density dependence of the pion decay constant as it enters in the chiral s-wave potential (22). We have combined this s-wave potential, treated in the local density approximation, with the time-honoured non-local p-wave potential which systematically reproduces the binding energies and widths of the higher-lying pionic atom states previously measured in stopped π^- experiments [11, 17].

Our results for pionic ^{207}Pb are shown in Fig 2. The points denoted "f_π" are obtained using the vacuum value of the pion decay constant in $T^{(-)}$ as it enters the s-wave optical potential (22), (23). We have used $\mathrm{Re}B_0 = 0$ in our "standard" set and Im $B_0 \simeq 0.06 m_\pi^{-4}$ following [11, 17]. The neutron radius is taken about 3% larger than the proton radius in the local densities $\rho_n^{(r)}$ and $\rho_p^{(r)}$. Clearly, the "vacuum f_π" scenario is quite far off

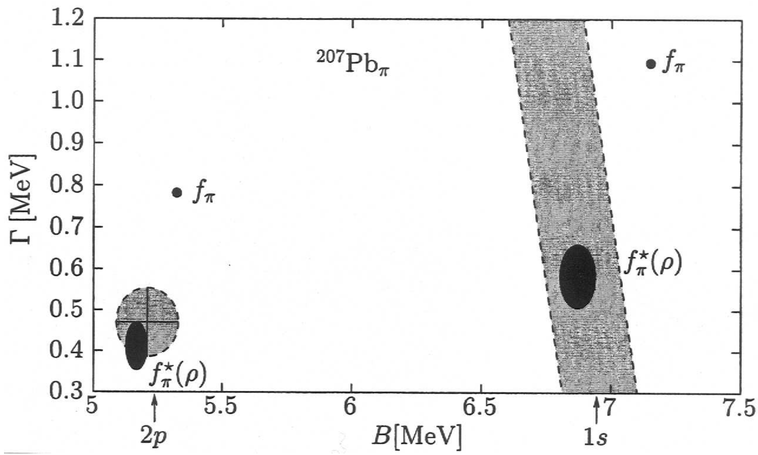

Fig. 2. Binding energy B and width Γ of $1s$ and $2p$ pionic atom states in ^{207} Pb. Points (f_π) are obtained [16] using the chiral s-wave optical potential (22, 23) with vacuum pion decay constant ($f_\pi = 92.4$ MeV) and $\mathrm{Re}B_0 = 0$. Dark ellipses (f^*_π) are results [16] when replacing f_π by the in-medium decay constant (25) with $\sigma_N = (45 \pm 8)$ MeV. Light shaded areas: empirical range of B, Γ from Ref. [12].

the lightly shaded areas which give the range of $1s$ and $2p$ binding energies and widths as deduced from the ^{208}Pb $(d,^3$ He$)^{207}$Pb$_\pi$ data [12]. The missing s-wave repulsion can of course be generated by simply adjusting $\mathrm{Re}B_0$. This would require a large negative value, $\mathrm{Re}B_0 \simeq -0.07 m_\pi^{-4}$, which would be at odds with theoretical many-body calculations [18], though within large uncertainties [19]. It would also be at odds with the π^--deuteron scattering length as mentioned earlier. On the other hand, replacing the vacuum pion decay constant in $T^{(-)}$ by $f^*_\pi(\rho)$ as given by Eq. (25), with ρ treated as local density distribution, the missing repulsion in the s-wave optical potential is easily supplied. The calculated results are shown by the dark ellipses in Fig 2. The data are now well reproduced using $\mathrm{Re}B_0 = 0$.

Our predictions for the case of pionic ^{205}Pb are shown in Fig 3 together with recent (preliminary) data [20]. This is a particularly interesting example because here the $1s$ state has been found as a well isolated peak in the $(d,^3$ He) spectrum.

4. Concluding remarks and perspectives

The results in Section 3 clearly demonstrate that detailed high precision studies of deeply bound pionic atoms do provide strong additional constraints on the s-wave pion–nucleus optical potential, especially when these studies are carried systematically through isotopic chains of neutron-rich

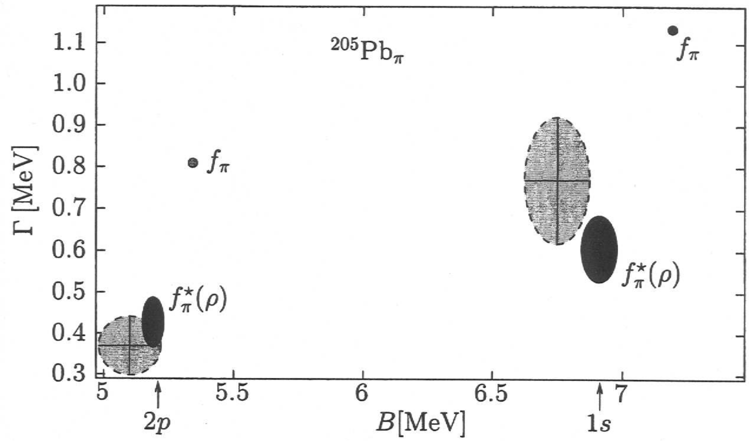

Fig. 3. Same as Fig 2, for pionic $1s$ and $2p$ states in ^{205} Pb. Theoretical predictions [16] are compared with preliminary data (light shaded ellipses) [20].

nuclei. We have taken the position here that the repulsion in the s-wave pion–nuclear interaction required to generate narrow $1s$ and $2p$ states, is naturally linked to the density dependence of the pion decay constant which in turn reflects the change of the QCD vacuum structure in dense matter.

We note that the isospin even and odd terms in the s-wave optical potential for a π^- are both repulsive, whereas for a π^+, the odd part changes sign and becomes attractive. For nuclei with $N/Z \simeq 1.5$ it turns out that the even and odd terms almost cancel, leaving a very small in-medium mass shift to the π^+. This appears to be compatible with the analysis of π^+ electroproduction on ^{3}He leading to ^{3}H in the final state [21].

We have omitted a variety of other interesting topics related to chiral dynamics in a nuclear medium, such as the splitting of K^+ and K^- masses which may have already been observed in high-energy heavy ion collisions at GSI. It is likely that the symmetry breaking pattern starting from the relatively weak π^+/π^- mass splitting in asymmetric nuclear matter proceeds via the $K/\bar{K}$ system to the splitting of D^+ and D^- masses in matter which offers the unique opportunity to investigate the in-medium dynamics of a *single* light quark attached to a heavy spectator quark.

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