Parton fluxes and virtual pions in heavy nuclei

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Abstract. The partonic flux originated from a heavy nucleus is not the mere sum of the fluxes coming from the individual nucleons. There are various effects that give rise to modifications. One of these effects is here investigated, i.e. the presence of a cloud of virtual pions co-moving with the nucleus. It is found that the contribution of these virtual particles to the total parton flux should be rather small even if one includes the contribution of resonances.

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1 Introduction

In collisions at very high energy involving nuclei, in particular, heavy nuclei, there is an interplay between high-energy phenomenology, which is described usually starting from different versions of parton models and the nuclear dynamics: the presence of many sources for the parton flux may induce relevant effects because they cannot be considered independent. These effects are usually studied directly at the partonic level and have therefore, at least in principle, their theoretical counterpart in QCD, even though it must be supplemented with phenomenological inputs to be able to produce answers at non-perturbative level, see e.g. [1]. Some of these investigations are moreover specifically devoted to the heavy-nuclei structure [2].

It is, however, clear that the nuclear interactions are not the direct expression of QCD, rather they are described by an intermediate dynamics which arises from the presence of a pion field interacting with the nucleons, so that the difference of the parton flux generated by a nucleus of number \( A \), from the flux produced by \( A \) independent nucleons arises also from the presence of interacting pions and, obviously, from the parton structure of the pions themselves.

The aim of this note is to analyze this particular aspect of the parton structure of heavy nuclei by means of a two-level model: a nuclear level in which the components are hadrons and a deeper partonic level in which the components are quarks and gluons. The model is not intended to reproduce details of the nuclear dynamics, but only to take into account the fundamental aspects of the pion-nucleon interaction as sources of the parton-flux modification. In a simplified dynamics, where there is only a quantized field of pions [3] linearly coupled to a source given by the nucleon, the effect is found to be very small. This result is easily recognized to be due to the particular pseudoscalar coupling, so other forms of coupling need to be investigated and this will be done by introducing the lightest of the mesonic resonances. These dynamical models are then applied to the system of many nucleons in order to see how the presence of many sources modifies the meson population with respect to the simple sum of the populations pertaining to the individual nucleons. There could be a problem of partial double counting of the same degree of freedom, but what is attempted to calculate is not the total pionic population but only that part of it which differ from the mere sum of the pionic populations produced independently by the single nucleons. The problem of the pion population of the nucleus, as is seen from scattering processes, has been already considered, see e.g. [4]. In order to be definite a precise dynamical model has been used, the one proposed in the book of J.D. Walecka [5].

This is done in the next section and the overall effect is found to be small. For this reason the detailed partonic structure of the pion, which has been investigated since long time [6], has not been used in detail. The effect looked for could have some relevance in the situation where many partonic interactions are relevant like in the so-called semi-hard processes in nucleon-nucleus or nucleus-nucleus interactions, where the total sub-energy of the single collision is high enough to allow a perturbative treatment, but the fractional momentum is small, so the partonic population involved is large and the rescattering probability is large. Due to the roughness of the model no evolution with \( Q^2 \) in the parton distribution is calculated. At the end a short estimate of the effects of Fermi motion...
and of correlations is presented, and within the limits of the model it is possible to state that both effects must be really small for heavy nuclei.

2 Virtual pions

2.1 Overview of the single-nucleon models

The models of Lewis, Oppenheimer and Wouthuysen [2] and of Bloch and Nordseick [7] are well known but a short overview of them is needed for the future discussions, in order to fix the notation and also because some modifications are introduced. (For a sketch of a different treatment, see the appendix). Both foresee a heavy spinor (the nucleon), which moves relativistically and interacts with lighter bosons. The nucleons cannot be created nor destroyed, whereas the bosons are emitted and absorbed, so the nucleon is treated in first quantization, the bosons by means of a second quantized field. The interaction has the standard form

\[ \mathcal{L}_\pi = ig_\pi \bar{\psi} \gamma_5 \vec{F} \cdot \vec{\phi} \bar{\psi} \]

for direct pion emission.

As anticipated further interactions are considered: this will be done, following the suggestion by Walecka, by introducing resonances, the \(\rho\) and the \(\omega\) spin-one mesons and moreover a scalar-isoscalar \(\sigma\) or \(f_0\). So we have three further interaction terms [5], respectively for the \(\omega\) and the \(\sigma\):

\[ \mathcal{L}_\rho = g_\rho \bar{\psi} \gamma_\mu \vec{U} \cdot \vec{\psi}, \quad \partial_\mu \vec{U} = 0; \]
\[ \mathcal{L}_\omega = g_\omega \bar{\psi} \gamma_\mu \vec{U} \cdot \vec{\psi}, \quad \partial_\mu \vec{U} = 0; \]
\[ \mathcal{L}_\sigma = g_\sigma \bar{\psi} \chi \bar{\psi}. \]

The dynamics is limited in two ways: there in no emission or absorption of nucleons, the emitted particles carry a small amount of momentum, so that the velocity of the nucleon is not perturbed in a sizable way. All these limitations fit well with the aim of describing the emission and absorption of pions inside the nucleus. So the total Hamiltonian is

\[ \mathcal{H}_T = -i\alpha_j \partial_j + \beta M + \sum_n \mathcal{H}_{0,n} \]
\[ -ig_\pi \beta \gamma_5 \vec{F} \cdot \vec{\phi} - g_\rho \beta \chi - g_\omega \beta \rho \cdot [\alpha_j \vec{U}_j - \vec{U}_0], \]

where \(\mathcal{H}_{0,n}\) are the free Hamiltonians of the pions, of the \(\rho\)-mesons and of the \(\omega\)-mesons, i.e. \(n = \pi, \sigma, \rho, \omega\). For the masses and the energies the notations are: \(M, E\) for the nucleon, \(m, \epsilon_n, \epsilon_{\pi}\) for the mesons. The wave function of the nucleon is written as \(\Psi(r) = \psi \sigma^f \tau^i\), where

\[ \psi = \sqrt{\frac{E + M}{2E}} \left[ \frac{\sigma^f \tau^i}{E + M} \right] \]

1. The arrow \(\vec{a}\) is used to denote isovectors, the corresponding components take an index \(u\), the boldface \(\bf{a}\) denotes spatial three-vectors, and \(j\) is the corresponding index; the nucleons will bear the indices \(i, \ell\), the different mesons the index \(n\). The normalization \(\xi^* \xi = 1\) gives \(\psi^* \psi = 1\). The Dirac matrices are evaluated between plane waves of given velocity with the known results:

\[ (\alpha_j) = \epsilon_j, \quad (\beta) = 1/\gamma, \quad (\gamma \epsilon_\pi) = \sigma \cdot \vec{k}/M \gamma. \]

The first non-zero result for \((\gamma \epsilon_\pi)\) is obtained at the first order in the momentum transfer; \(\vec{k} \equiv (k_\perp, k_\parallel/\gamma)\) and the directions transverse and parallel are defined with respect to the velocity \(v\). The vector fields are decomposed taking into account the subsidiary condition \(\partial_\mu U^n = 0\), so that only three polarizations enter [8]. At this point there are still matrix structures, the spin matrices \(\sigma_j\) and the isospin matrices \(\tau_n\), they make the subsequent calculation much more complicated than in Q.E.D. [7]. They will be substituted by numerical vectors \(s_j, t_u\) corresponding again to some mean value over the nucleon state; the normalization is such that \(s_j\) is twice the spin and \(t_u\) twice the isospin of a nucleon. In so doing, the original Hamiltonian becomes

\[ \mathcal{H}_T = -i \vec{v} \cdot \partial + M/\gamma + \sum_n \int \{ \epsilon_n, \epsilon_n^f \} a^\dagger_n(a_n(k)) d^3k. \]

The sources \(J_n\) have the explicit form:

\[ J_n = \frac{i g_\pi}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\epsilon_n \gamma}} \vec{s} \cdot \vec{k} e^{i \vec{k} \cdot \vec{r}}, \]
\[ J_s = \frac{g_\rho}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\epsilon_n \gamma}} e^{i \vec{k} \cdot \vec{r}}, \]
\[ J_\rho_n = \frac{g_\rho}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\epsilon_n \gamma}} V_j^f e^{i \vec{k} \cdot \vec{r}}, \]
\[ J_\omega_n = \frac{g_\omega}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\epsilon_n \gamma}} V_j^f e^{i \vec{k} \cdot \vec{r}}. \]

The vectorial coefficients \(V_j\) are given by

\[ V_j = e_j \cdot v, \quad \vec{k} \cdot e_j = 0, \quad \text{for } j = 1, 2, \]
\[ V_3 = \frac{e \cdot k}{m|k|} - \frac{|k|}{m}. \]

It is known that the total Hamiltonian is brought to a diagonal form by means of the transformation

\[ \mathcal{U}^\dagger \mathcal{H}_T \mathcal{U} = b_n, \]

where

\[ \mathcal{U} = \exp \sum_n \int \{ a_n F_n - a_n^\dagger F_n^* \} d^3k \]

with

\[ F_n = f_n e^{i \vec{k} \cdot \vec{r}} = \frac{J_n}{\epsilon_n - v \cdot k_n}. \]

In this way, in the lowest level of \(\mathcal{H}_T\), characterized by \(b_n(k)|\psi\rangle = 0\), the meson population is given by

\[ \langle \psi | a_n^\dagger a_n | \psi \rangle = |F_n|^2 = |f_n|^2. \]

The total population is then obtained by summing over the internal (spin and isospin) quantum numbers.

The populations are independent of \(r\), they depend, however, on \(v\).