

MESONIC AND BINDING CONTRIBUTIONS TO THE NUCLEAR DRELL-YAN PROCESS

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Abstract

We have evaluated the Drell-Yan cross section in nuclei paying special attention to the meson cloud contribution from pion and ρ -meson, for which an accurate calculation using the meson nuclear spectral functions is used. Similarly, the nucleonic contribution is evaluated in terms of a relativistic nucleon spectral function. Fair agreement with experiment is found for different nuclei and the results show a sizeable contribution from the renormalized meson cloud. In order to reproduce the experiment a novel element is introduced, consisting of a gradual energy loss of the incoming proton in its pass through the nucleus which produces a strong A dependence at x_1 large.

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The Drell-Yan process (DY) in nuclei in which a quark (antiquark) from a proton beam and an antiquark (quark) from the constituents of the nuclear target fuse to produce a $\mu^+\mu^-$ pair, has been advocated as a complement of deep inelastic scattering (DIS) in nuclei in order to investigate the reasons why the quarks in a nuclear medium (even if they are integrated into other effective degrees of freedom, like nucleons and mesons) behave differently than when they are in isolated nucleons [1, 2, 3, 4, 5, 6, 7]. Comprehensive reviews of this process and its relation to the EMC effect can be seen in [8, 9, 10, 11].

Since in the DY process one needs an antiquark for the reaction, there was the hope that, compared to the DIS, the nuclear DY would be more sensitive to modifications of the pion cloud in the nucleus, one of the mechanisms originally suggested to explain the EMC effect [8, 9, 10, 11].

The calculations in ref. [2] show that the region of large x_1 and small x_2 in the Drell-Yan process is enhanced due to medium modifications of the pion cloud, while the region of small x_1 would simply reproduce the same results as the EMC effect. In ref. [5] the pionic excess leads to a sharp increase of the DY ratio for large x_1 and $x_2 < 0.5$. Similarly an increase of that ratio with increasing values of $x_1 \cdot x_2$ is predicted in [3] due to mesonic effects.

The experimental data, however, does not show the expected enhancement [12] and recent papers try to justify the small role of the pion cloud in this process and consequently in the EMC effect [13]. On the other hand, the mesonic contribution in the EMC effect was revived recently in [14], where a new evaluation was carried out in terms of the pion spectral function, avoiding the static approximation (use of pion excess number) and demanding that the pion propagator satisfies rigorous analytical properties.

The mesonic effects of ref. [14] showed up as an enhancement of the ratio F_{2A}/F_{2D} in the region $0.1 < x < 0.5$, in qualitative agreement with previous findings. However, the contribution from the renormalized pion cloud was found smaller than claimed in the past [9, 10]. At the same, time the renormalized ρ -meson cloud in the medium was shown to be as relevant as the one of the pion, and even more important in some kinematical regions. Together with the use of an accurate nucleon spectral function, including relativistic effects, the analysis of [14] showed a good agreement with experiment in the EMC region for a wide range of nuclei.

In the present work we would like to extend the ideas of ref. [14] to the Drell-Yan process and establish the role of the meson cloud in that reaction. The results show an important role of the meson cloud in the DY process but, unlike in earlier works based upon the meson cloud, we can now find agreement with experiment. In ref. [7] the effects of the pion cloud in the DY process are also reanalyzed demanding similar requirements as those of ref. [14]. The effects of the ρ -meson cloud are not considered there.

There is another important novelty in our work with respect to former works in the DY process. Since the protons of the beam are strongly interacting particles, they will have strong collisions with the nucleons much more often than they will eventually produce the $\mu^+\mu^-$ pair. We shall take this into

consideration here. We shall work outside the shadowing region for the variable x_2 in order to avoid the typical coherent phenomena of this region. Outside this shadowing region the different parts of the nuclear volume will contribute incoherently to the cross section. The incoming proton collides strongly with the nucleons, but the proton keeps travelling and there is no loss of flux. However, in any of these collisions the proton of the beam will lose a certain amount of energy. This will change the variable x_1 and consequently the contribution to the cross section. Let us see how this is implemented:

The cross section for $pN \rightarrow \mu^+\mu^-X$ is given by

$$d^2\sigma(pN \rightarrow \mu^+\mu^-X) = \frac{4\pi\alpha^2 K}{9q^2} \sum_a e_a^2 [q_a(x_1)\bar{q}_a(x_2) + \bar{q}_a(x_1)q_a(x_2)] dx_1 dx_2 \quad (1)$$

where the sum is over the flavours of the quarks and x_1, x_2 refer to the beam and target nucleons respectively. The variables x_1, x_2 indicate the fraction of the momentum carried by the quark (antiquark) of the beam and target nucleons which fuse to create the virtual photon that leads to the $\mu^+\mu^-$ pair of momentum q . In an invariant form we have

$$x_1 = 2\frac{q \cdot p_2}{s} \quad ; \quad x_2 = 2\frac{q \cdot p_1}{s} \quad ; \quad s x_1 x_2 = q^2 \quad (2)$$

where p_1, p_2 are the fourmomenta of the beam and target nucleons and $s = (p_1 + p_2)^2$. In the frame where the target nucleon is at rest we can write:

$$x_1 = \frac{q^0}{E_1} \quad ; \quad x_2 = \frac{q^0 - q^3}{M} \quad (3)$$

where M is the nucleon mass and E_1 the energy of the incoming nucleon. The axis 3 is chosen along the direction of the beam.

On the other hand we shall have a distribution of p_2^0, p_2^3 in the nucleus given by the nuclear spectral function. This will give rise to a variable x_{2N} which in the Bjorken limit can be written in terms of the static x_2 variable as

$$x_{2N} = 2\frac{q \cdot p_1}{s_N} = x_2 \frac{M}{p_2^0 - p_2^3} \quad (4)$$

which is the same relationship of x_N to x in DIS [14].

The amount of energy that the nucleon of the beam loses in one collision with the nucleons in the nucleus is difficult to quantize. On the one hand some energy is transferred to the target and many new particles can be created at the high energies which we are discussing, $E_1 \simeq 800$ GeV. On the other hand some times the nucleon can get excited to some resonant state in the strong collision and the resonance may behave like a nucleon with respect to the electromagnetic $\mu^+\mu^-$ process. Furthermore, given the scale of time in which the process occurs, the $\mu^+\mu^-$ production could take place before the asymptotic final state from the strong collision materializes. All this tells about

the difficulties in determining the equivalent energy loss for one collision. For this reason we do not attempt to evaluate this magnitude, but use the same DY experiment to fix it. We assume that in each collision, occurring with a probability $\sigma_{NN}\rho dl$, we lose a fraction β of the energy. This fraction does not need to be the same for all energies, but we will assume it to be constant in the energy range where we move. We take the same function for computations in different nuclei. The value of β is taken such as to reproduce one experimental point in one nucleus, which will be described later on. Thus we have:

$$\begin{aligned}\frac{dE_1}{dl} &= -\sigma_{NN}\rho\beta E_1 \\ E_1(\vec{r}) &= E_{1in} \exp[-\beta\sigma_{NN} \int_{-\infty}^z \rho(\vec{b}, z') dz'] \\ x_1(\vec{r}) &= x_1 \exp[\beta\sigma_{NN} \int_{-\infty}^z \rho(\vec{b}, z') dz']\end{aligned}\quad (5)$$

where σ_{NN} is the NN total cross section, $\sigma_{NN} = 40$ mb, $\rho(\vec{r})$ the nuclear density and \vec{b} the impact parameter. We have assumed that the nucleon loses energy but keeps moving in the forward direction, in the spirit of the eikonal approximation. In proton elastic collisions with the nucleon ($\sigma_{el} \simeq \sigma_{Tot}/6$) the amount of energy lost is negligible because the cross section is very forward peaked [15]. So, strictly speaking we should use σ_{in} instead of σ_{NN} , but since what matters is the product $\beta \cdot \sigma$ and β is fitted to the data, we can keep using the formalism of eq. (5).

With this prescription and the formalism of ref. [14] we obtain

$$\begin{aligned}\frac{d^2\sigma^{(N)}(pA \rightarrow \mu^+\mu^-X)}{dx_1dx_2} &= \frac{4\pi\alpha^2K}{9q^2} 4 \int d^3r \sum_a e_a^2 [q_a(x_1(\vec{r})) \\ &\times \int \frac{d^3p}{(2\pi)^3} \frac{M}{E(\vec{p})} \int_{-\infty}^{\mu} dp^0 S_h(p^0, p) \bar{q}_a(x_{2N}) + \bar{q}_a(x_1(\vec{r})) \\ &\times \int \frac{d^3p}{(2\pi)^3} \frac{M}{E(\vec{p})} \int_{-\infty}^{\mu} dp^0 S_h(p^0, p) q_a(x_{2N})] \theta(x_{2N}) \theta(1-x_{2N}) \theta(1-x_1(\vec{r}))\end{aligned}\quad (6)$$

where $S_h(p^0, p)$ is the hole spectral function for the nucleon in the nucleus.

For the pion cloud contribution we have, following ref. [14]

$$\begin{aligned}\frac{d^2\sigma^{(\pi)}(pA \rightarrow \mu^+\mu^-X)}{dx_1dx_2} &= \frac{4\pi\alpha^2K}{9q^2} (-6) \int d^3r \\ &\sum_a e_a^2 [q_a(x_1(\vec{r})) \int \frac{d^4p}{(2\pi)^4} \theta(p^0) \delta Im D_\pi(q) 2M \bar{q}_a(x_{2\pi}) \\ &+ \bar{q}_a(x_1(\vec{r})) \int \frac{d^4p}{(2\pi)^4} \theta(p^0) \delta Im D_\pi(q) 2M q_a(x_{2\pi})] \\ &\theta(x_{2\pi} - x_2) \theta(1-x_{2\pi}) \theta(1-x_1(\vec{r}))\end{aligned}\quad (7)$$

The ρ contribution is obtained from eq. (7) changing $D_\pi(q) \rightarrow D_\rho(q)$, $x_{2\pi} \rightarrow x_{2\rho}$ and the factor (-6) by (-12) .

The magnitude $\delta D(p)$ is given, following ref. [14], by

$$\delta D(p) = D(p) - D_0(p) - \left. \frac{\partial D}{\partial \rho} \right|_{\rho=0} \cdot \rho \quad (8)$$

with $D(p)$ standing for the π or ρ propagator in the medium and $D_0(p)$ the corresponding free one. Furthermore

$$x_{2\pi} = x_2 \frac{M}{p_2^3 - p_2^0} \quad (9)$$

p_2 referring now to the pion momenta. The expression of $x_{2\pi}$ compared to the one of x_{2N} in eq. (4) has an apparent minus sign which comes because the variable p for pions has opposite direction as for nucleons for convenience. The variable $x_{2\rho}$ is given by eq. (9) but now p_2 stands for the ρ momentum. The extra factor 2 appearing for the ρ meson contribution in the cross section, with respect to the expression for pions, accounts for the two transverse polarizations of the ρ which couple to the nucleons in our approach [14].

In eq. (6) the q, \bar{q} with arguments $x_1(\vec{r})$ refer to the proton beam and $q_a(x_{2N})$ refer to the nucleon target (average over proton and neutron).

In eq. (7), instead, the q, \bar{q} with argument $x_{2\pi}(x_{2\rho})$ refer to the quark distributions of the pion (ρ meson, this latter one is taken the same as for pions [16, 17]). The quark distributions of nucleons and pions are taken from refs. [18] and [19] respectively. For the meson propagator and nucleon spectral function we use the same input as in the study of the EMC effect in ref. [14].

A small technical change is made here. The form factor accompanying the coupling of the mesons to the $ph, \Delta h$ components is taken here of the monopole type as in [14], with the same values of Λ , but we take it static (dependent on three momentum only) to avoid unrealistic contributions from regions in the integration where one is close to the poles of the form factor. We should note that these form factors are constructed in the study of the NN interaction where $q^2 < 0$ [20]. While this does not practically modify the results in DIS, it produces some reduction of the results here. Particularly, the meson contribution to the DY in one nucleon, which can be obtained by substituting $\delta D(p)$ by the last term in eq. (8), is more strongly reduced and reaches reasonable numbers.

Next we show some results. We have evaluated the results for $R = 2d\sigma_A/Ad\sigma_D$ corresponding to the experiment of [12]. We have evaluated the results in two different ways. First we take a value of x_2 and evaluate a weighed cross section multiplying the cross sections at different values of x_1 by the experimental acceptances [21]. As indicated in [12], we include in the average only the regions $4 \text{ GeV} \leq \sqrt{Q^2} \leq 9 \text{ GeV}$ and $\sqrt{Q^2} \geq 11 \text{ GeV}$, in order to avoid the regions of quarkonium resonances. Another calculation is done by assuming $x_1 - x_2 = 0.26$ (the average value in [12]) in order to approximately extrapolate the results to the unmeasured region. We do not evaluate $d\sigma$ for the deuteron in our local density approach and hence we divide by the average of the cross

section on the proton and the neutron. For values of $x_2 < 0.6$ this has errors of less than 2 % as found in the study of DIS [22, 23].

We fix the fraction of energy loss in order to obtain one point of the spectrum of ^{56}Fe , at $x_2 = 0.15$. The resulting fraction is $\beta = 0.035$, or 3.5% energy loss. This is a small quantity but it plays a role as we shall see.

In figs. 1,2,3 we show the results for ^{12}C , ^{40}Ca and ^{56}Fe . The experimental results are roughly reproduced in the three nuclei. We notice the relative importance of using the weighed cross section instead of the one using the average value of $x_1 - x_2$, although the differences are smaller than 10%. The latter cross sections show a small A dependence which is different than the one of the weighed cross sections. The ratios obtained from the weighed cross sections are of the order of unity, as approximately shown by the experiment, although the ^{12}C data seems to be larger than unity around $x_2 \simeq 0.25$. The results using the $x_1 - x_2$ average also show a ratio bigger than unity for ^{12}C in this region.

It is interesting to show the effects of the meson cloud. In figs. 1,2,3 we have separated the nucleonic contribution from the one of the pion and ρ -meson clouds. As we can see, the contribution of the pion cloud is a bit larger than the ρ one. The role of the energy loss can be seen in fig. 1. At small values of x_2 , taking $x_1 - x_2 = 0.26$, it reduces the cross section in about 10 %. At larger values of x_2 the reduction is much bigger. This stronger reduction is easily understood since, due to the energy loss, x_1 increases from its original value $x_2 + 0.26$ and eventually becomes bigger than unity where there is no strength for the cross section.

It is interesting to see, however what are our predictions in two limiting cases, x_1 small and x_1 large. In fig. 4 we show the results for the three nuclei for $x_1 = 0.01$ and $x_1 = 0.7$. For $x_1 = 0.01$ we obtain values of R very similar to those obtained for the EMC effect, as already noticed in [2]. The A dependence in this case is very weak. However, for $x_1 = 0.7$ the dependence of R on x_2 is quite different for all three nuclei. It is easy to understand this different behaviour: for x_1 very small, the energy degradation increases x_1 by a certain fraction but it is still very small and this energy loss induces small changes (smaller than 10 %). On the contrary, when x_1 is close to unity, a fractional increase of x_1 brings it closer to one where there is no strength. This induces large reductions in the cross sections, which are larger in heavier nuclei, hence the A dependence seen in the cross sections.

Although our pionic effects are qualitatively similar to those in earlier works, the energy loss changes appreciably the spectra with respect to former predictions [2, 5].

The issue of the energy loss in hadronic processes is attracting some attention. Early estimates of quark energy losses based on the uncertainty principle pointed towards very small energy losses of quarks propagating through the nucleus [24]. More accurate estimates, yet not free of uncertainties as claimed by the authors, are done in [25].

These estimates are lower than 3.5% energy loss per collision. More recent

evaluations [26] distinguish between quarks created in the nuclear medium and incoming quarks. For the latter case an energy loss proportional to the energy is obtained which can be translated in our language as about 2.5% energy loss per collision. On the other hand, the energy loss per collision counted asymptotically, is certainly larger than 3.5% in order to reproduce the broad energy distributions of the experimental cross sections [27] (see also comments in ref. [28] and particularly fig. 8). These experimental facts can be reconciled recalling our arguments that in some cases hadronic resonances are excited, carrying the energy of the incoming proton, which propagate through the nucleus and can produce $\mu^+\mu^-$ similarly as nucleons. However, in pure hadronic reactions these resonances would decay outside the nucleus into a nucleon and mostly pions and the nucleon would have less energy than the incoming one.

It is also interesting to quote that in a recent paper [29] looking at the propagation of J/ψ in nuclei a solution was favoured implying both absorption of J/ψ through the nucleus and energy loss of the beam. This loss was equivalent to 3.6% per collision, although some small trade-offs could be made between the absorption and the energy loss.

Given the relevance of the mesonic components and the energy loss in the DY process, we should worry about questions of selfconsistency since the pion structure functions are determined from analysis of DY processes in which no energy loss is assumed. A look at fig. 4 would tell us that for values of x_1 small, where the energy loss played a small role, one has not much to worry, but for large values of x_1 such things could be more relevant. This would add certain uncertainties to our predictions for $x_1 = 0.7$ in fig. 4.

In any case the study done here clearly shows that the issue of the energy loss is an important one and experimental efforts should be devoted to clarify it. Fig. 4, even with accepted uncertainties, also shows that the DY process at large values of x_1 is the relevant place to look at.

In summary, our results show a sensitivity of R to the meson cloud renormalization in nuclei, more important than in the EMC effect. The values of R around unity in our interpretation are not a signal of the lack of mesonic effect, but the simultaneous effect of the mesonic cloud and the progressive energy degradation of the nucleon beam through the nucleus. The latter has the effect of reducing the cross section since x_1 increases and $q(x_1)$ ($\bar{q}(x_1)$) decreases. The present interpretation of the DY nuclear effect has as a consequence a stronger nuclear dependence than in DIS, where the ratio is practically constant for $A > 7$. Here we have found that for x_1 close to unity the ratio R is rather dependent on A . It would be very interesting to have other DY experiments done with a larger range of values of x_1, x_2 and better precision in order to be able to test the novel consequences that the present interpretation of the process provides.

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Figure Captions

- Figure 1: The ratio $2d\sigma_A/Ad\sigma_D$ for the Drell-Yan process on ^{12}C . Experimental points from ref. [12]. Long-dashed line: nucleonic contribution. Dashed-dotted line: plus pionic contribution. Solid line: plus ρ -meson contribution (full calculation). The former curves use the experimental acceptances and cuts. Short dashed line: full calculation using $x_1 - x_2 = 0.26$. Dotted line: full calculation with $x_1 - x_2 = 0.26$ omitting the energy loss.

- Figure 2: Same as fig. 1 for ^{40}Ca .

- Figure 3: Same as fig. 1 for ^{56}Fe .

- Figure 4: Three upper curves around small x_2 : values of R for fixed value $x_1 = 0.01$ for ^{12}C , ^{40}Ca , ^{56}Fe (from up down), as a function of x_2 . The lower curves around small x_2 : values of R for fixed value $x_1 = 0.7$ for ^{12}C , ^{40}Ca , ^{56}Fe (from up down), as a function of x_2 .







