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Comments and Addenda

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A dependence of valence- and sea-quark contributions to Drell-Yan processes

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The valence-quark contribution to the Drell-Yan process can be isolated by data with particle and antiparticle beams. It is suggested that the A dependence of the sea-quark contributions could be larger than the linear A increase of the valence contributions. The implications of such a sea having quarks with large momenta are discussed.

The quark structure function of the charged pion has recently been determined from the data on dimuon production in $\pi^{\pm}N$ collisions[1-3] in the context of the Drell-Yan model.[4] In this model the production of large-mass dimuons in hadron-hadron collisions proceeds by the fusion of a quark in one hadron with the antiquark in the other hadron. The nuclear targets provide antiquarks only through their sea, whereas in projectiles such as $\pi^{\pm}$, K^-, and $\bar{p}$ antiquarks enter as the dominant valence components. Sea-quark distributions are not precisely known. The sea needed to fit the data on pN scattering is in general agreement with the sea distributions measured by the spacelike probes in the deep-inelastic lepton scattering experiments.[2,3] On the whole, however, reliable measurements of sea-quark distributions do not exist except possibly through the Drell-Yan formula itself. In practice several assumptions are made in the parametrization of the sea, such as whether or not it is SU(3) symmetric, whether the sea includes charm and other heavier quark flavors, etc.[5] Such assumptions could then markedly influence analyses of dimuon production in pN collisions. This is in contrast to the analyses of deep-inelastic scattering wherein the sea contributions are usually correction terms. In order to determine the pion structure function $u^{\pi}(x)$ it would be desirable to isolate the valence contributions independently of the sea distribution functions.

Another important ingredient in the Drell-Yan analyses is the A dependence of the cross section where A is the mass number of the target nucleus. As the rates of dimuon production are quite small, experiments are done on targets of heavy nuclei (to gain a factor of about one hundred) and the "per nucleon" cross section is extracted by dividing by A^{α}. This procedure is extremely sensitive to the precise value of α; for instance, a change of 10% in the value of α can change the overall normalization by as much as 100% (i.e., to a ratio $K \sim 2$). In fact the experimental value of $\alpha = 1.12 \pm 0.05$ of the Chicago-Illinois-Princeton (CIP) group[6] differs from the value obtained by the CERN NA-3 group[2] $\alpha = 1.03 \pm 0.03$. We suggest that this discrepancy arises because (i) the A^{α} factor does not adequately represent the true A dependence and (ii) the two groups are measuring A dependences of different observables. The reason for this may well be that the sea of the target nucleus is not a simple sum of the seas of the constituent nucleons, but is in fact larger.

In the following we list the combinations of the valence distributions which contribute to the Drell-Yan process. The valence contributions are isolated by taking the difference between particle beams and antiparticle beams on a given target. The only assumption is that the sea distribution of a specific flavor is the same as that of the corresponding antiflavor. Motivated partly by the discrepancy between the different experimental values of α in $\pi N \to \mu^+\mu^- + \cdots$ we suggest that the A dependence of the sea contributions is stronger than the A dependence of the valence contributions. The assumed A dependence of the sea of the nu-

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cleus provides a qualitative explanation to the rise of α in high-p_T events in inclusive hadron production.

SEPARATION OF VALENCE CONTRIBUTIONS

We start with the Drell-Yan formula[4] for the collision of a hadron beam particle B on a target nucleon N,

$$BN \to \mu^+\mu^- + \text{hadrons}\,,$$

$$\frac{d^2\sigma^{BN}}{dx_1 dx_2} = \left(\frac{4\pi\alpha^2}{3s}\right)\frac{1}{3}\frac{1}{x_1^2 x_2^2} H^{BN}(x_1,x_2)\,, \tag{1}$$

$$H^{BN}(x_1,x_2) \equiv \sum_f Q_f^2 [\bar{f}^B(x_1) f^N(x_2) + f^B(x_1)\bar{f}^N(x_2)]\,. \tag{2}$$

The notation is standard. The kinematic relations of the Bjorken variables x_1 and x_2 to the Feynman variable x_F and the dimuon mass m are

$$x_1 - x_2 = x_F\,, \qquad x_1 x_2 = \frac{m^2}{s}\,. \tag{3}$$

The structure function $f(x)$ represents the parton distribution of the quark flavor f [$\bar{f}(x)$ for antiquarks]. The functions $f(x)$ are normalized such that $(1/x)f^B(x)$ is the probability of finding a quark of flavor f in hadron B with a fractional longitudinal momentum x. Instead of dealing with the double differential cross section we shall deal with the equivalent function H^{BN}.

For the case of a beam consisting of antiparticles $\bar{B}$ we have

$$H^{\bar{B}N}(x_1,x_2) = \sum_f Q_f^2[\bar{f}^{\bar{B}}(x_1) f^N(x_2) + f^{\bar{B}}(x_1)\bar{f}^N(x_2)] = \sum_f Q_f^2[f^B(x_1) f^N(x_2) + \bar{f}^B(x_1)\bar{f}^N(x_2)]\,, \tag{4}$$

where we have used invariance under charge conjugation. The difference between antiparticle $\bar{B}$ and particle B beams is then

$$\Delta H^{BN}(x_1,x_2) \equiv H^{\bar{B}N}(x_1,x_2) - H^{BN}(x_1,x_2) \tag{5}$$

$$= \sum_f Q_f^2 [\bar{f}^B(x_1) - f^B(x_1)][\bar{f}^N(x_2) - f^N(x_2)]\,. \tag{6}$$

Assuming that each structure function can be separated into a valence (v) part and a sea (s) part,

$$f(x) = f_v(x) + f_s(x)\,, \tag{7}$$

we have

$$\Delta H^{BN} = \sum_f Q_f^2 [(\bar{f}_v^B - f_v^B) + (\bar{f}_s^B - f_s^B)] \times [(\bar{f}_v^N - f_v^N) + (\bar{f}_s^N - f_s^N)]\,, \tag{8}$$

where the f and $\bar{f}$ in the first (second) brackets are functions of x_1 (x_2).

The difference of the cross sections will involve only the valence contributions if we assume

$$f_s^h(x) = \bar{f}_s^h(x)\,, \tag{9}$$

i.e., in a hadron h the sea-quark distribution function of a given flavor f is the same as the distribution function of the antiquark corresponding to that flavor. This assumption is weaker than assumptions usually made such as SU(3), SU(4) for the sea. For instance, the sea distributions in the proton $u_s^p, d_s^p, s_s^p, c_s^p, b_s^p, \ldots$ could be unrelated to each other and still satisfy the above assumption. On the other hand, it is possible that $u_s^p(x)$ may be slightly less than $\bar{u}_s^p(x)$ for moderate values of x because of the Pauli principle operating between the valence quarks $u_v^p(x)$ and the sea quarks $u_s^p(x)$. In the following we shall assume the validity of Eq. (9) and thus obtain from Eqs. (7), (8), and (9):

$$\Delta H^{BN}(x_1,x_2) = \sum_f Q_f^2 [f_v^B(x_1) - \bar{f}_v^B(x_1)] f_v^N(x_2)\,. \tag{10}$$

Thus the sea contributions of both the beam and target nucleon drop out and only the valence contributions remain in the difference.[7]

For a nucleus target, following the original parton model hypothesis, we simply add the contributions coming from the scattering on the *valence* quarks of each nucleon in the nucleus. Thus Eq. (10) implies that for a target nucleus of Z protons and N neutrons

$$\Delta H^{BA} = \tfrac{1}{9}[4(u_v^B - \bar{u}_v^B)(Zu_v^p + Nd_v^p) + (d_v^B - \bar{d}_v^B)(Zd_v^p + Nu_v^p)]\,. \tag{11}$$

We shall list below the above difference for the cases of experimental interest:

$$\frac{d^2\sigma^{\bar{B}A}}{dx_1dx_2} - \frac{d^2\sigma^{BA}}{dx_1dx_2} = \frac{4}{9}\frac{\pi\alpha^2}{sx_1^2x_2^2}\Delta H^{BA}(x_1,x_2)\,, \tag{12}$$

where

$$\Delta H^{pA} = \tfrac{1}{9}\{4u_v^p(x_1)[Zu_v^p(x_2) + Nd_v^p(x_2)] + d_v^p(x_1)[Zd_v^p(x_2) + Nu_v^p(x_2)]\}\,, \tag{13}$$

$$\Delta H^{\pi^+A} = \tfrac{1}{9}u_v^{\pi^+}(x_1)[(4Z - N)u_v^p(x_2) + (4N - Z)d_v^p(x_2)]\,, \tag{14}$$

$$\Delta H^{K^+A} = \tfrac{4}{9}u_v^{K^+}(x_1)[Zu_v^p(x_2) + Nd_v^p(x_2)]\,. \tag{15}$$

The functional dependence on x_1 and x_2 is factored in Eqs. (14) and (15). Therefore in order to completely determine the valence structure function of $\pi^{\pm}$, for example, it is sufficient to have accurate data on one nuclear target. The condition

$$\int_0^1 u_v^{\pi^+}(x)\,\frac{dx}{x}=1 \qquad (16)$$

fixes the normalization. Data with π^+ and π^- beams on more than one target can be used to check the consistency of our procedure. The problem of determining $u_v^{K^+}$ is analogous to the case of $\pi^{\pm}$. It is, however, not possible to determine the strange valence-quark distribution in the K meson $\bar{s}_v^{K^+}$ without additional assumptions such as SU(3) symmetry, etc.

We emphasize that this procedure of extracting structure functions is free of detailed assumptions about the sea, and independent of the deviations of the measured α from unity (which are due to the sea contributions). The functions u_v^p and d_v^p obtained by using pN and $\bar{p}N$ data in Eq. (13) refer to timelike momentum transfers and should be checked against those functions extracted from deep-inelastic lepton scattering on hadrons in the spacelike region. Moreover, the quantum-chromodynamic corrections would not affect the above procedure because although the structure functions acquire $\log Q^2$ dependence, due to the assumption of Eq. (9) the cancellation of the sea contributions goes through. It may, however, be noted that the $\log Q^2$ dependence observed by the CERN-Dortmund-Heidelberg-Saclay group in νN and $\bar{\nu}N$ experiments turn out for the process $\pi N\rightarrow\mu^+\mu^-+\cdots$ (Ref. 2) to be negligible (~5%) in the Q^2 range 16–80 $(\text{GeV})^2$.

A DEPENDENCE OF THE SEA

We substitute Eq. (7) in Eq. (2) and obtain for a target nucleon N

$$H^{BN}(x_1,x_2)=\sum_f Q_f^2[(\bar{f}_v^B+\bar{f}_s^B)f_v^N+(\bar{f}_v^B f_s^N+f_v^B\bar{f}_s^N)+(\bar{f}_s^B f_s^N+f_s^B\bar{f}_s^N)]\,. \qquad (17)$$

In going to the case of a nucleus A, we notice that the factor in the first parentheses involves only valence densities of the target and hence falls under the case discussed already (according to which the individual nucleon contributions are to be added leading to $\alpha=1$). The second and third terms are of the form (valence × sea) and (sea × sea) and thus involve the sea quarks of the nucleus. If we speculate that these terms acquire the dependence A^{α_s}, where α_s is greater than unity we have for a target nucleus,

$$H^{BA}\sim A[v(x_1)v(x_2)+s(x_1)v(x_2)]+A^{\alpha_s}[v(x_1)s(x_2)+s(x_1)s(x_2)]\,, \qquad (18)$$

which may be compared with the customary parametrization

$$H^{BA}\sim A^{\alpha}[v(x_1)v(x_2)+s(x_1)v(x_2)+v(x_1)s(x_2)+s(x_1)s(x_2)]\,. \qquad (19)$$

Our empirical reasons for expecting $\alpha_s>1$ are as follows: (i) At small positive values of $x_F=x_1-x_2$, the dominant term in the above sum comes from $v(x_1)v(x_2)$ which according to Eq. (18) is multiplied by A. For larger values of x_F the important term $v(x_1)s(x_2)$ is associated with A^{α_s}. Thus as we go to larger values of x_F we expect to see a gentle increase in the effective value of α according to Eq. (19). In fact an increase of α of about 15% with $x_F=0.5\rightarrow0.9$ is not ruled out in the CIP data using π^- beam on C, Cu, and W targets.[6] However, more accurate data on α are required to establish this trend. (ii) The average value of α measured in the CIP experiment[6] for dimuon mass $m>4$ GeV is

$$\alpha=1.12\pm0.05\,, \qquad (20)$$

whereas the value determined by the CERN NA-3 group[2] using π^+ and π^- beams on H and Pt is

$$\alpha=1.03\pm0.03\,. \qquad (21)$$

The CIP value uses the parametrization of Eq. (19). The NA-3 group isolates the valence contribution by the difference method using Eq. (14) and hence their A dependence according to Eq. (18) must be governed by the first term, leading to $\alpha=1$. Thus there is no contradiction between the two measurements.[8]

It may be remarked that the foregoing considerations have little effect on the determination of the shape of the pion structure function $u_v^{\pi}(x_1)$, but change the overall normalization by a factor $\sim A^{0.12}$.

Some comments about the sea of the nucleus may be in order. The $q\bar{q}$ sea of the nucleus consists of two parts: a part intrinsic to each constituent nucleon arising from vacuum fluctuations and the other part originating from the mutual interactions among the nucleons. The meson exchanges which are essentially $q\bar{q}$ exchanges between pairs of nucleons augment the sea. A high-Q^2 probe therefore sees the nucleus as being composed of nearly stationary nucleons, each with its three valence quarks and the intrinsic sea, embedded in a fluid of $q\bar{q}$ pairs generated by mutual interactions.[9] The intrinsic parts of the nucleons add up to a contribution which goes as A, whereas the part of the sea arising from nuclear binding could vary as A^{α_s}, where α_s is slightly larger than unity.[10]

The effect of the Pauli principle is generally to make some of the sea quarks (antiquarks) occupy higher-momentum states. If the number of sea quarks grows like A^{α_s}, then the corresponding Fermi momentum $k_F \sim A^{(\alpha_s-1)/3}$, which makes the sea "harder." In terms of the usual parametrization of the sea components $u_{sea}(x) \sim (1-x)^\beta$, etc., we therefore expect β to decrease[11] as we increase the mass number A. This hardening of the sea component in momentum space, if correct, has some interesting consequences. As some of the sea quarks and antiquarks have larger momenta it will simulate situations which imply the presence of large Fermi motions of nucleons in nuclei. The presence of energetic quarks in the sea will enhance the probability for the so-called "hard scattering." This could then provide an explanation for the observed increase of α with increasing p_T in hadron-nucleus collisions.[12]

In conclusion, valence-quark contributions can be isolated by the method of taking differences. It will be important to check whether they grow linearly with A in all the three sets of reactions $BA \to \mu^+\mu^- + \cdots$, where $B = \pi^\pm, K^\pm, p^{(-)}$. If the contributions of the sea to the Drell-Yan process increase faster than linearly with A, then it may be possible to understand why the two πN experiments[2,6] observe two apparently different values for the α parameter. It would be extremely interesting to check whether this is a valid explanation by finding that the α defined through Eq. (19) increases with the Feynman variable x_F.

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[7] It may be noted that since the gluon distributions in the particles B and $\bar{B}$ are the same, $G^B(x) = G^{\bar{B}}(x)$, the quantum-chromodynamic corrections arising from diagrams involving arbitrary number of gluons but at most one q or $\bar{q}$ in the initial channel of the subprocess (e.g., the Compton process $gq^{(-)} \to \gamma q^{(-)}$) also cancel in the differnece (10).

[8] In this connection it may be noted that the α in pN collisions reported by the Columbia–Fermilab–Stony Brook group [Ref. 5 and D. F. Kaplan *et al.*, Phys. Rev. Lett. 40, 435 (1978)], although averaged over a wide range of dimuon masses, refers to values of x_F in a narrow range around $x_F = 0$ due to their limited acceptance. For instance, their experiment has the range $x_F = -0.15$ to $+0.075$ at 400 GeV/c. It would be important to check whether there is an increase in α by going to larger x_F.

[9] The picture is somewhat akin to the Bloch's theory of the free electrons in metals, where the positive ion cores are stationary and the conduction electrons move about almost freely throughout the volume of the metal.

[10] According to the saturation of nuclear binding, the effective number of nucleon pairs in interaction is roughly linear in A but not a higher power of A. However, the binding energy of a nucleus depends on the time-averaged picture, whereas the sea that is seen by a high-Q^2 probe is an instantaneous state of affairs inside the nucleus. For long times the distant interactions of a nucleon in the nucleus are shielded by the $q\bar{q}$ pairs surrounding it. I thank Professor K. W. McVoy for a discussion on this point.

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