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Obtaining information on short range correlations from inclusive electron scattering

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In view of recent data from the Thomas Jefferson National Accelerator Facility (JLab) on inclusive electron scattering off nuclei at high momentum transfer ($Q^2 \gtrsim 1$ GeV^2) and their current analysis, it is shown that, if the scaling variable is properly chosen, the analysis in terms of scaling functions can provide useful information on short-range correlations (SRC). This is demonstrated by introducing a new relativistic scaling variable that incorporates the momentum dependence of the excitation energy of the $(A-1)$ system, with the resulting scaling function being closely related to the longitudinal momentum distributions.

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Obtaining information on short-range correlations (SRC) in nuclei is a primary goal of modern nuclear physics [1]. Interest in SRC stems not only from the necessity to firmly establish the limits of validity of the standard model of nuclei but also from the impact that the knowledge of the detailed mechanism of SRC would have in understanding the role played by quark degrees of freedom in hadronic matter and the properties of the latter in dense configurations [2]. Recently, evidence of SRC has been provided by new experimental data on inclusive $[A(e,e')X]$ [3,4] and exclusive $[A(e,e'pN)X$ and $A(p,pN)X]$ lepton and hadron scattering off nuclei at high momentum transfer ($Q^2 \gtrsim 1$ GeV^2) (see Ref. [5] and references therein quoted). In inclusive scattering the observation of a scaling behavior of the ratio of the cross section on heavy nuclei to that on the deuteron [3], for values of the Bjorken scaling variable $1.4 \lesssim x_{\rm Bj} \lesssim 2$, and to that on ^{3}He [4], for $2 \lesssim x_{\rm Bj} \lesssim 3$, has been interpreted as evidence that the electron probes two- and three-nucleon correlations in complex nuclei similar to the ones occurring in two- and three-nucleon systems [6]. It should be pointed out, however, that whereas exclusive processes can directly access the relative and center-of-mass motions of a correlated pair in a nucleus [7], obtaining information on these quantities from inclusive scattering is, in principle, more difficult. Various approaches based on scaling concepts have therefore been proposed, going from the already mentioned scaling behavior of the cross section ratio plotted versus $x_{\rm Bj}$ to the scaling behavior of the ratio of the nuclear to the nucleon cross sections plotted versus proper scaling variables; among the latter, a process that has been most investigated in the past is the so-called Y-scaling, for it is believed that this may represent a powerful tool to extract the high-momentum part of the nucleon momentum distribution, which is governed by SRC [8,9]. It is the aim of this Rapid Communication to critically reanalyze the concept of Y-scaling, mainly because of (i) the lack of a general consensus about the usefulness of such a concept and (ii) a strong renewal of interest in Y-scaling owing to recent experimental data on $A(e,e')X$ reactions from the Thomas Jefferson National Accelerator Facility (JLab) [10,11]. We will show that the analysis of inclusive scattering in terms of proper Y-scaling variables could indeed provide useful information on SRC; to this end, following the suggestion of Refs. [12–14] a new approach to Y-scaling and its usefulness will be illustrated in detail. Let us consider a virtual photon of high momentum impinging on a nucleus A (with mass M_A) and knocking out, in a quasielastic process, a nucleon N (with mass m_N) having momentum $k \equiv |\mathbf{k}|$ and removal energy E. The latter is defined as the energy necessary to remove the nucleon from A, leaving the residual nucleus $(A-1)$ (with mass M_{A-1}) with intrinsic excitation energy E^*_{A-1} (i.e., $E = m_N + M_{A-1} - M_A + E^*_{A-1} = E_{\rm min} + E^*_{A-1}$). In the plane wave impulse approximation (PWIA) and using the instant form of dynamics, the quasielastic cross section reads as follows:

$$\sigma_2^A(q,\nu) \equiv \frac{d^2\sigma(q,\nu)}{d\Omega_2\, d\nu} = \sum_{N=1}^{A} \int dE\, d\mathbf{k} P_N^A(k,E)\sigma_{en} \times (q,\nu,\mathbf{k},E)\,\delta(\nu + M_A - E_N - E_{A-1}), \quad (1)$$

with energy conservation ($M^*_{A-1} = M_{A-1} + E^*_{A-1}$)

$$\nu + M_A = \sqrt{m_N^2 + (\mathbf{k}+\mathbf{q})^2} + \sqrt{M^{*^2}_{A-1} + \mathbf{k}^2} \quad (2)$$

and momentum conservation $\mathbf{q} = \mathbf{p} + \mathbf{p}_{A-1}$. Here $\nu = \epsilon_1 - \epsilon_2$ and $\mathbf{q} = \mathbf{k}_1 - \mathbf{k}_2$ are the energy and three-momentum transfers ($Q^2 = q^2 - \nu^2 = 4\epsilon_1\,\epsilon_2 \sin^2\frac{\theta}{2}$ with $q \equiv |\mathbf{q}|$), σ_{en} is the elastic electron cross section off a moving off-shell nucleon with momentum $k \equiv |\mathbf{k}|$ and removal energy E, and $P_N^A(k,E)$ is the spectral function (normalized to one) of nucleon N (i.e., the joint probability to have a nucleon with momentum k and removal energy E); eventually, $\mathbf{p}$ and $\mathbf{P}_{A-1}$ are the momenta of the undetected struck nucleon and the final $(A-1)$ system. Considering, for ease of presentation, isoscalar nuclei, one has $P_N^A(k,E) = P_p^A(k,E) = P_n^A(k,E) \equiv P^A(k,E) = P_0^A(k,E) + P_1^A(k,E)$, where $P_0^A(k,E) = (1/A)\sum_{\alpha\in F} A_\alpha n_\alpha(k)\,\delta(E-\epsilon_\alpha)$ is the (trivial) shell-model part [with A_α denoting the occupation number of the single-particle state α with removal energy ϵ_α and momentum distribution $n_\alpha(k)$] and P_1 is the (interesting) part generated by NN correlations. The spectral function is linked to the momentum distributions by the momentum sum rule $n^A(k) = \int P^A(k,E)\,dE = \int P_0^A(k,E)\,dE + \int P_1^A(k,E)\,dE = n_0^A(k) + n_1^A(k)$. It has been shown [8] that at high values of momentum transfer, after integrating over the direction of **k**, Eq. (1) can be written, to a good approximation, as follows:

$$\sigma_2^A(q,\nu) \simeq \left\{[Zs_{ep}(q,\nu,k,E) + Ns_{en}(q,\nu,k,E)]\frac{E_p}{q}\right\}_{(k_{\min},E_{\min})} \times F^A(q,\nu), \quad (3)$$

where s_{en} is the electron-nucleon cross section integrated over the polar angle and $F^A(q,\nu)$ is the nuclear structure function

$$F^A(q,\nu) = 2\pi\int_{E_{\min}}^{E_{\max}(q,\nu)} dE \int_{k_{\min}(q,\nu,E)}^{k_{\max}(q,\nu,E)} k\,dk P^A(k,E). \quad (4)$$

Equation (3) is obtained by eliminating the δ function by integrating over $\cos\alpha = (\mathbf{k}\cdot\mathbf{q}/kq)$, with the limits of integration resulting from the condition $-1 \leqslant \cos\alpha \leqslant 1$. We can now introduce a scaling variable $Y = Y(q,\nu)$, which is only required to be a function of q and ν (and any arbitrary constant) so that, regardless of the specific form of Y, the cross section and structure function can be expressed, without loss of generality, in terms of the two independent variables q and $Y = Y(q,\nu)$, rather than the canonical q and ν. Correspondingly, a scaling function $F^A(q,Y)$ is introduced; this is nothing but Eq. (4) with ν replaced everywhere by Y; if, under certain conditions, $F^A(q,Y) \to F^A(Y)$, Y-scaling is said to occur and, depending on the physical meaning of Y and $F^A(Y)$, information on nucleons in nuclei could be obtained. To simplify our analysis, let us consider high values of the momentum transfer, when $E_{\max}(q,Y)$ and $k_{\max}(q,Y,E)$ become so large that, because of the rapid falloff of $P^A(k,E)$, they can be replaced by ∞ (although in actual calculations we use the correct values of these quantities); in this case, the q and ν dependence of the scaling function is governed only by $k_{\min}(q,Y,E)$, and it is trivial to show that, by adding and subtracting a proper term, the scaling function can be cast in the following general form:

$$F^A(q,Y) = 2\pi\int_{E_{\min}}^{\infty} dE \int_{k_{\min}(q,Y,E)}^{\infty} k\,dk P^A(k,E) = f^A(Y) - B^A(q,Y), \quad (5)$$

where $f^A(Y) = 2\pi\int_{|Y|}^{\infty} k\,dk n^A(k)$ represents the *longitudinal momentum distribution*, and

$$B^A(q,Y) = 2\pi\int_{E_{\min}}^{\infty} dE \int_{|Y|}^{k_{\min}(q,Y,E)} k\,dk P_1^A(k,E) \quad (6)$$

is the *binding correction* [8], which, through $k_{\min}(q,Y,E)$, is governed by the continuum energy spectrum of the final $(A-1)$ system, unlike $f^A(Y)$, which is integrated over all excited states of $(A-1)$. The quantities $f^A(Y)$ and $n^A(k)$ are linked by the relation $n^A(k) = -[df^A(Y)/dY]/[2\pi Y]$, $k = |Y|$, so that if $f^A(Y)$ could be extracted from the experimental data, $n^A(k)$ could be determined. Unfortunately, the presence of $B^A \neq 0$ such an extraction is hindered by depending upon the difference between Y and $k_{\min}$ and therefore upon the definition of the former. The binding correction is absent only in the deuteron, since $E = E_{\min} = \text{constant} = 2.22$ MeV, so that $Y = k_{\min}(q,\nu,E_{\min})$, $B^D(q,Y) = 0$, and $F^D(q,Y) = f^D(Y)$. The final state interaction (FSI) of the struck nucleon invalidates the PWIA, but, in spite of that, an approach was developed in the past to reduce the effects from both the binding corrections and FSI [8]; the approach is based upon the widely used relativistic scaling variable $Y = y$ [8–11,15], which is obtained by setting, in the energy conservation equation [Eq. (2)], $k = y$, $\mathbf{k}\cdot\mathbf{q}/kq = 1$, and, most importantly, $E_{A-1}^* = 0$; thus y represents the *minimum longitudinal momentum of a nucleon having the minimum value of the removal energy* $E = E_{\min}$. In the asymptotic limit ($q \to \infty$), Eq. (5) scales in y and becomes the *asymptotic scaling function* $F^A(y) = f^A(y) - B^A(y)$, that is, Eq. (5) with Y and $k_{\min}(q,Y,E)$ replaced by y and $k_{\min}^{\infty}(y,E)$, respectively (scaling in this variable also occurring within a relativistic description of the deuteron [15]). Unfortunately, owing to the presence of $B^A(y)$, $F^A(y)$ is not related to a momentum distribution so that, in principle, the experimental longitudinal momentum distribution $f_{\text{ex}}^A(y)$ and, consequently, $n_{\text{ex}}^A(k)$, cannot be extracted from the data. Let us briefly recall how this problem was addressed in Ref. [8]. The experimental scaling function $F_{\text{ex}}^A(q,Y) = \sigma_{2,\text{ex}}^A(q,Y)/\{[Zs_{ep}(q,\nu,k,E) + Ns_{en}(q,\nu,k,E)](E_p/q)\}_{(k_{\min},E_{\min})}$ exhibits, when $Y = y$, a strong q dependence owing to the FSI and binding effects and differs from the asymptotic scaling function $F_{\text{ex}}^A(y)$. The latter, however, has been obtained in Ref. [8] by extrapolating to $q \to \infty$ the available values of $F_{\text{ex}}^A(q,y)$, on the basis that FSI can be represented as a power series in $1/q$ and dies out at large q^2, a conclusion that has been reached by various authors (see, e.g., Ref. [16]). The experimental longitudinal momentum distribution $f_{\text{ex}}^A(y)$ has thereby been obtained by adding to $F_{\text{ex}}^A(y)$ the binding correction $B^A(y)$ evaluated theoretically, and $n_{\text{ex}}^A(k)$ has been obtained by $n^A(k) = d[F^A(y) + B^A(y)/dy]/[2\pi y]$, $k = |y|$. Such a procedure affects the final results in terms of large errors on the extracted momentum distributions, particularly at large values of k; in spite of these errors, the extracted momentum distributions at $k \gtrsim 1.5$–$2\ \text{fm}^{-1}$ turned out to be larger by orders of magnitude from the prediction of mean-field approaches and in qualitative agreement with realistic many-body calculations that include SRC. To make the extraction of $f_{\text{ex}}^A(y)$ as independent as possible from theoretical binding corrections, in Ref. [12] another scaling variable $Y = y_{\text{CW}}$ has been introduced; this scaling variable incorporates relevant physical dynamical effects left out in the definition of y. To readily understand the physical meaning of the new scaling variable, let us consider the asymptotic limit of $k_{\min}(y,q,E)$ for a large nucleus [i.e., $k_{\min}^{\infty}(y,E) = |y - (E - E_{\min})|$]; it can be seen that only when $E = E_{\min}$ does $k_{\min}(y,E) = |y|$, in which case $B^A = 0$ and $F^A(y) = f^A(y)$; this holds only for the deuteron, whereas for a complex nucleus $E_{A-1}^* \neq 0$ and $E \geqslant E_{\min}$, so $B^A(y) \neq 0$, and $F^A(y) \neq f^A(y)$. It is therefore the dependence of $k_{\min}$ on E_{A-1}^* that gives rise to the binding effect [i.e., to the relation $F^A(y) \neq f^A(y)$]. This is an unavoidable defect of the usual approach to Y-scaling, based on the scaling variable y; in fact, the longitudinal momentum is very different for weakly bound shell-model nucleons ($E_{A-1}^* \sim 0$–20 MeV) and strongly bound correlated nucleons ($E_{A-1}^* \sim 50$–200 MeV), and at large values of $|y|$ the scaling function is not related to the longitudinal momentum of strongly bound correlated nucleons, whose contributions almost entirely exhaust the behavior of the scaling function. As stressed in Refs. [12–14], to establish a global link between experimental data and longitudinal momentum components, one has to conceive a scaling variable *that could equally well represent longitudinal momenta of both weakly bound and strongly bound nucleons so that the binding correction could be minimized*. This can be achieved by adopting a scaling variable that properly includes the momentum dependence of the average excitation energy of $(A-1)$ generated by correlations, namely,

$$\langle E^*_{A-1}(k)\rangle = \frac{1}{n^A(k)}\int P_1^A(k, E^*_{A-1})E^*_{A-1} d\,E^*_{A-1}, \qquad (7)$$

where $E^*_{A-1} = E - E^{(2)}_{\rm thr}$, $E^{(2)}_{\rm thr} = M_{A-2} + 2m_N - M_A$ being the threshold energy for two-particle emission. We have calculated the quantity in Eq. (7) using a realistic spectral function for nuclear matter and ^{3}He. The results are presented in Fig. 1, where they are compared with the prediction of the spectral function of the few-nucleon correlation (FNC) model of Ref. [19], according to which

$$E^*_{A-1}(\mathbf{k}, \mathbf{K}_{\rm CM}) = \frac{A-2}{A-1}\frac{1}{2m_N}\left[\mathbf{k} - \frac{A-1}{A-2}\mathbf{K}_{\rm CM}\right]^2, \qquad (8)$$

where $\mathbf{K}_{\rm CM}$ is the CM momentum of a correlated pair. In view of the very good agreement between the FNC model and the exact many-body results for nuclear matter and ^{3}He, we used the former to calculate $\langle E^*_{A-1}(k)\rangle$ for nuclei with $3 < A < \infty$. The values shown in Fig. 1 can be interpolated by

$$\langle E^*_{A-1}(k)\rangle = \frac{A-2}{A-1}T_N + b_A - c_A|\mathbf{k}|, \qquad (9)$$

where $T_N = (\sqrt{m_N^2 + k^2} - m_N)$ and b_A and c_A result from the CM motion of the pair ($b_{\rm NM} = 37.3$ MeV, $c_{\rm NM} = 0.04$ and $b_3 = -2.94$ MeV, $c_3 = -0.03$). Placing in Eq. (2) $k = y_{\rm CW}$, $\frac{\mathbf{k}\cdot\mathbf{q}}{kq} = 1$, and $M^*_{A-1} = M_{A-1} + \langle E^*_{A-1}(k)\rangle - \langle E_{\rm gr}\rangle$, we obtain a fourth-order equation for the new scaling variable $y_{\rm CW}$, which, in contrast to previous work [12–14], has been solved exactly; this, together with the relativistic extension of the definition of the mean excitation energy, is necessary to extend $y_{\rm CW}$ to high values. Note that the value of $\langle E_{\rm gr}\rangle$, fixed by the Koltun sum rule (see Refs. [12–14]), has been added to Eq. (9) to counterbalance the effects of $\langle E^*_{A-1}\rangle$ at low $y_{\rm CW}$. For a large nucleus and not too large values of $y_{\rm CW}$, one has

$$y_{\rm CW} = -\frac{\tilde q}{2} + \frac{\nu_A}{2W_A}\sqrt{W_A^2 - 4m_N^2}. \qquad (10)$$

Here, $\nu_A = \nu + \tilde M_D$, $\tilde M_D = 2m_N - E^{(2)}_{\rm th} - b_A + \langle E_{\rm gr}\rangle$, $\tilde q = q + c_A\nu_A$, and $W_A^2 = {\nu_A}^2 - \mathbf{q}^2 = \tilde M_D^2 + 2\nu\tilde M_D - Q^2$. For the deuteron $y_{\rm CW} = y = -q/2 + (\nu_D/2W_D)\sqrt{W_D^2 - 4m_N^2}$ with $\nu_D = \nu + M_D$ and invariant mass $W_D^2 = {\nu_D}^2 - \mathbf{q}^2 = {M_D}^2 + 2\nu M_D - Q^2$; for small values of $y_{\rm CW}$, such that $\frac{A-2}{A-1}(\sqrt{y^2_{\rm CW} + m_N^2} - m_N) + b_A - c_A|y_{\rm CW}| \ll \langle E_{\rm gr}\rangle$, the variable y, representing the longitudinal momentum of a weakly bound nucleon, is recovered. Therefore $y_{\rm CW}$ effectively takes into account the k dependence of E^*_{A-1}, both at low and high values of $y_{\rm CW}$, and interpolates between the correlation and the single-particle regions; it can be interpreted as the *minimum longitudinal momentum of a nucleon that, at high values of $y_{\rm CW}$, has removal energy $\langle E^*_{A-1}\rangle$ and is partner of a correlated two-nucleon pair with effective mass $\tilde M_D$.*

Let us now illustrate the merits of $y_{\rm CW}$-scaling and its practical usefulness. The main merit is that, because of the definition of $y_{\rm CW}$, binding effects play a minor role, as clearly illustrated in Fig. 2; practically, $k_{\rm min}(q,\nu,E) \simeq |y_{\rm CW}|$ and $B^A(q, y_{\rm CW}) \simeq 0$, with two relevant consequences: (i) to a large extent $F^A(q, y_{\rm CW}) \simeq f^A(y_{\rm CW})$ [cf. Eq. (5)] and (ii) as a result of (i), one would expect that at high values of $y_{\rm CW}$, $F^A(q, y_{\rm CW})$ will behave in the same way in the deuteron and in complex nuclei, since $n^A(k) \simeq C_A n^D(k)$ and, accordingly, $f^A(y_{\rm CW}) \simeq C_A f^D(y_{\rm CW})$; at low values of $y_{\rm CW}$, in contrast, $F^A(q, y_{\rm CW})$ should exhibit an A dependence generated by the different

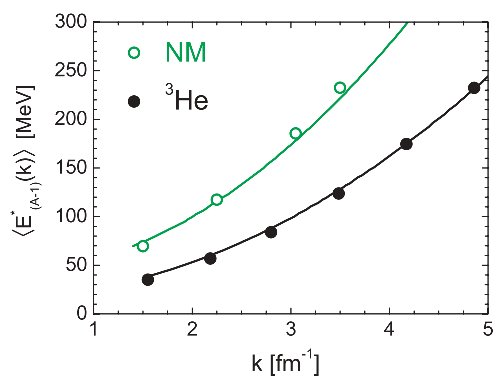

FIG. 1. (Color online) The average value of $E^*_{A-1}(k)$ [Eq. (7)] calculated for nuclear matter with the spectral function of Ref. [17] (open dots) and for ^{3}He with the spectral function from the Pisa wave functions [18] (full dots). The full lines are obtained with the spectral function of the few-nucleon correlation model of Ref. [19].

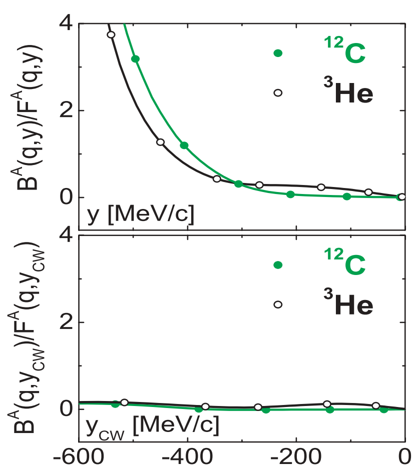

FIG. 2. (Color online) The ratio of the binding correction [Eq. (6)] to the scaling function [Eq. (5)] for ^{3}He (open dots) and ^{12}C (full dots) calculated with the scaling variable y, which does not contain any effective excitation energy from SRC (upper panel), and with the variable $y_{\rm CW}$, which takes into account SRC effects by Eq. (9) (lower panel).

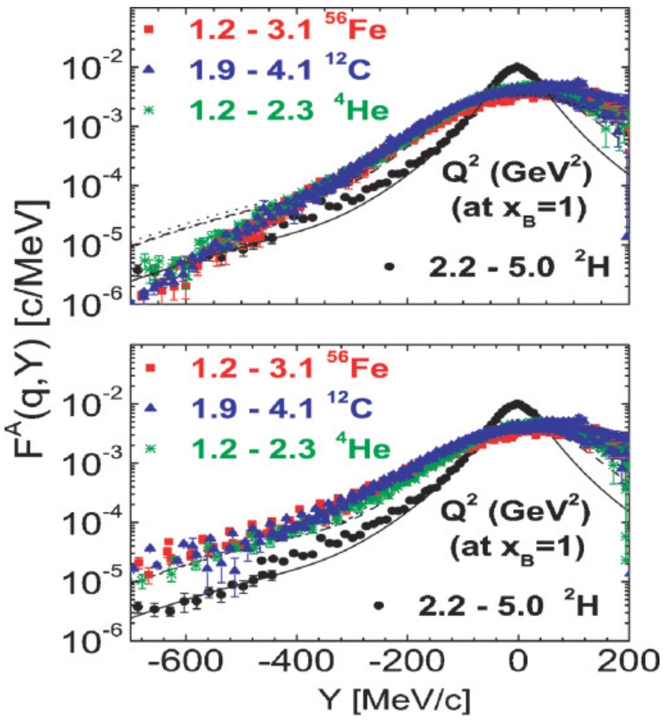

FIG. 3. (Color online) The experimental scaling function (symbols) for ^{4}He, ^{12}C, and ^{56}Fe obtained from the experimental data of Refs. [10,21]. The upper panel shows $F^A(q, Y = y)$ and the lower panel $F^A(q, Y = y_{\rm CW})$. The full, long-dashed, dashed, and dotted curves represent the longitudinal momentum distributions $f^A(Y) = 2\pi \int_{|Y|}^{\infty} n^A(k)kdk$ for ^{2}H, ^{4}He, ^{12}C, and ^{56}Fe, respectively, calculated with realistic wave functions.

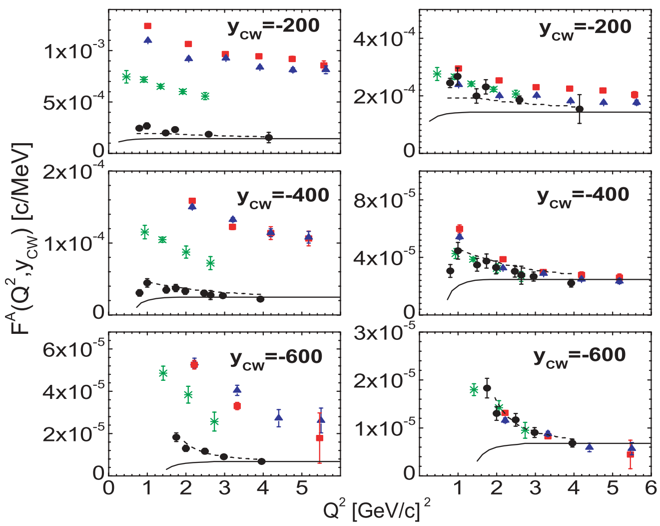

FIG. 4. (Color online) The scaling function $F^A(Q^2, y_{\rm CW})$ from the lower panel of Fig. 3 plotted vsQ^2 at fixed values of $y_{\rm CW}$ (^{4}He, asterisks; ^{12}C, triangles; ^{56}Fe, squares). In the right panel the data for ^{4}He, ^{12}C, and ^{56}Fe have been divided by the constants $C_4 = 2.7$, $C_{12} = 4.0$, and $C_{56} = 4.6$, respectively. The theoretical curves refer to ^{2}H and represent the PWIA results (full) and the results that include the FSI (dashed), both obtained with the AV18 interaction [20]. Scaling variables are in MeV/c.

asymptotic behavior of the nuclear wave functions in configuration space. This is fully confirmed in Fig. 3, which, moreover, also shows that whereas $F^A(q, y)$ scales to a quantity that strongly differs from the longitudinal momentum distribution, $F^A(q, y_{\rm CW})$ scales exactly to $f^A(y_{\rm CW})$. This is even better demonstrated in Fig. 4, where the effects of FSI are also illustrated. The left panel shows that (i) scaling is violated and approached from the top (which is clear signature of the breaking down of the PWIA, which has to approach scaling from the bottom [8]) and (ii) the Q^2 dependence of the scaling violation appears to be the same for the deuteron and complex nuclei, a fact that has never been demonstrated before and represents, in our opinion, a relevant finding. To better validate point (ii), we have divided $F^A(Q^2, y_{\rm CW})$ by a constant C_A, such as to obtain $F^A(Q^2, y_{\rm CW})/C_A \simeq F^D(Q^2, y_{\rm CW})$. The results are shown in the right panel of Fig. 4; it can again be seen that not only at high values of $|y_{\rm CW}|$ do all scaling functions scale in A, but, more importantly, the constants C_A agree, within the statistical errors, with the theoretical predictions of Ref. [6], as well as with the experimental results on the ratio $R(x_{\rm Bj}, Q^2) = 2\sigma_2^A(x_{\rm Bj}, Q^2)/A\sigma_2^D(x_{\rm Bj}, Q^2)$ [3]. The main findings of our analysis can be summarized as follows: (i) at high values of $|y_{\rm CW}|(\gtrsim$200–300 MeV/c) the scaling function $F^A(Q^2, y_{\rm CW})$ scales to the one of the deuteron, with scaling constants C_A in qualitative agreement with theoretical predictions and other types of experimental analysis; this kind of A-scaling is entirely due to the scaling of the momentum distributions, $n^A(k) \simeq C_A n^D(k)$, at $k \gtrsim 1.5$–2 fm^{-1}, which can therefore be investigated by $y_{\rm CW}$-scaling analysis of inclusive data, owing to the direct link between the scaling function $F^A(Q^2, y_{\rm CW})$ and the longitudinal momentum distributions; (ii) the FSI has relevant effects on the scaling functions up to $Q^2 \simeq 4$–5 GeV^2 but, most importantly and surprisingly, it exhibits a similar Q^2 dependence in complex nuclei and in the deuteron; this has neither been observed nor theoretically predicted previously; in a forthcoming paper it will indeed be shown that the effects of the FSI on the momentum distribution of a correlated nucleon are similar in the deuteron and in a complex nucleus (for preliminary results see Ref. [22]).

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