\( \Xi^- \) and \( \Omega \) Distributions in Hadron-Nucleus Interactions\(^1\)

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ABSTRACT

Strange baryons have long been known to exhibit a leading particle effect. A recent comparison of \( \Xi^- \) production in \( \pi^- \), \( n \), and \( \Sigma^- \) interactions with nuclei show this effect clearly. These data are supplemented by earlier measurements of \( \Xi^- \) and \( \Omega \) production by a \( \Xi^- \) beam. We calculate the \( \Xi^- \) and \( \Omega \) \( x_F \) distributions and nuclear dependence in \( hA \) interactions using the intrinsic model.

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1 Introduction

Leading particle effects, flavor correlations between the final-state hadron and the projectile valence quarks, have long been observed in strange particle production. Although many experiments have recently focused on leading charm production [1, 2, 3, 4, 5, 6, 7, 8, 9, 10], the first data involved strange particles [11, 12, 13, 14, 15, 16]. With new data from the WA89 collaboration on Ξ⁻ production by π⁻, n, and Σ⁻ projectiles on nuclear targets [17], in addition to Ξ⁻ production data from Ξ⁻ beams [14], doubly strange hadron production can be studied as a function of the number of strange valence quarks in the projectile. We compare our model calculations to both the \( x_F \) distributions and the integrated \( A \) dependence reported by WA89 [17]. We also discuss Ξ⁻ and Ω production by the Ξ⁻ beam [14].

The WA89 collaboration used carbon, C, and copper, Cu, targets to study the \( A \) dependence of Ξ⁻(dss) production by π⁻(ud), n(udd), and Σ⁻(dds) beams [17]. The negative beams, π⁻ and Σ⁻, had an average momentum of 345 GeV with a 9% momentum spread. The neutron beam had a lower momentum with a larger spread than the negative beams—the average momentum was 260 GeV with a 15% variation. The detected Ξ⁻ was in the forward \( x_F \) region, \( x_F \geq 0.05 \), with low transverse momentum, \( p_T \leq 2.5 \) GeV. The data were parameterized in the form

\[
\frac{d\sigma}{dp_T^2dx_F} \propto (1-x_F)^a e^{-bp_T^2}. \tag{1}
\]

The pion and neutron results agree with the functional form of Eq. (1) over all \( x_F \). For the pion, \( a = 3.8 \pm 0.3 \) for C and 4.1 ± 0.3 for Cu while for the neutron \( a = 5.0 \pm 0.3 \) for C and 4.8 ± 0.3 for Cu. These results are consistent with expectations from spectator counting rules [18], \( d\sigma/dx_F \propto (1-x_F)^{2n_s-1} \). With an incident gluon, \( n_s = 2 \) for pions and 3 for neutrons, consistent with no leading particle effect for projectiles with zero strangeness. There is no significant \( A \) dependence of the exponent \( a \).

On the other hand, the Σ⁻ data cannot be fit to Eq. (1) for \( x_F < 0.4 \). In the large \( x_F \) region, \( a = 2.08 \pm 0.04 \) for C and 1.97 ± 0.04 for Cu. These results indicate a very hard \( x_F \) distribution, inconsistent with the counting rules even for a valence quark since \( n_s = 2 \) gives \( (1-x_F)^3 \). In addition, at \( x_F < 0.4 \), the distribution is independent of \( x_F \) for both targets. Thus these data show a strong leading particle effect since the Ξ⁻ has two valence quarks in common with the Σ⁻. The statistics are also sufficient for an observable \( A \) dependence in the fitted values of \( a \).

The integrated \( A \) dependence was also reported by WA89 [17]. The \( A \) dependence of the total cross section is often parameterized as

\[
\sigma_{pA} = \sigma_{pp}A^\alpha. \tag{2}
\]

The integrated \( \alpha \) for Σ⁻ production of Ξ⁻, \( \alpha = 0.679 \pm 0.011 \) [17], is in relatively good agreement with previous fits. However, the pion and neutron data show a closer-to-linear
A dependence, $\alpha = 0.891 \pm 0.034$ and $0.931 \pm 0.046$ respectively. WA89 attributes this difference to the fact that two $s\bar{s}$ pairs must be produced to make the final-state $\Xi^-$ and $s\bar{s}$ pair production would be suppressed relative to light $q\bar{q}$ production.

WA89 has also measured the dependence of $\alpha$ on $x_F$. This dependence, $\alpha(x_F)$, was previously reported for a wide range of hadron projectiles [19]. For non-strange hadrons and hadrons with a single strange quark, there is a common trend with $x_F$. At $x_F = 0$, $\alpha \approx 0.8$ and decreases to $\approx 0.5$ at large $x_F$, an overall decrease of $\sim A^{1/3}$ for $0 < x_F < 1$. The $\Xi^0$, the only doubly-strange hadron included in Ref. [19], is an exception. In $pA$ interactions, the $\Xi^0$ has a larger value of $\alpha$ at low $x_F$ [16]. A similar effect is observed for $\Xi^-$ production by WA89. Their measurements of $\alpha(x_F)$ for $\Xi^-$ from pion and neutron beams show that $\alpha \sim 1$ for $x_F \sim 0.05$, decreasing to $\alpha \sim 0.7$ at higher $x_F$. Thus the decrease of $\alpha$ with $x_F$ is also $A^{1/3}$ in this case although the actual values of $\alpha$ are larger than those for lighter hadrons [19]. However, for $\Sigma^-$-induced $\Xi^-$ production, $\alpha \sim 0.7$ almost independent of $x_F$.

The other data we consider are $\Xi^- \rightarrow \Xi^-, \Omega(sss)$ at 116 GeV, measured by Biagi et al. [14]. In this case, the final-state $\Xi^-$ $x_F$ distribution increases with $x_F$, as does the $\Omega$ $x_F$ distribution. This increase could be due in part to the use of an invariant parameterization [14],

$$E \frac{d\sigma}{dp^3} \propto (1 - x_F)^{a'} e^{-b'p_T^2},$$

which fits the $\Xi^-$ data at $x_F > 0.5$ but only approximately fits the $\Omega$ data in this limited region. The exponent $a'$ was fit in two $p_T^2$ intervals, $p_T^2 < 0.4$ GeV$^2$ and $0.4 < p_T^2 < 2.9$ GeV$^2$, yielding $a' = -0.45 \pm 0.02$ and $-0.18 \pm 0.03$ respectively. Between the most central measurement, $x_F = 0.15$, and the projectile fragmentation region, $x_F = 0.85$, the $\Xi^-$ cross section increases by a factor of $\sim 40$ in the low $p_T^2$ interval.

A comparison of these results with incident proton data [14], $pA \rightarrow \Xi^- X [11, 20]$ and $pA \rightarrow pX [21]$, showed that, at low $x_F$, the $\Xi^-$ production cross section is essentially independent of the projectile while, at high $x_F$, the $\Xi^-$ and $p$ scattering cross sections are similar. This behavior supports valence quark domination at high $x_F$. The structure of the $\Omega$ $x_F$ distribution is similar: it is of the same order of magnitude as $pA \rightarrow \Omega X$ [11] at low $x_F$ but is similar to singly strange baryon production by protons, $pA \rightarrow \Lambda X$ and $\Sigma^+ X$, [11, 20, 22] at high $x_F$.

Since only one target was used, $\alpha = 0.6$ was assumed in Eq. (2) to obtain the per nucleon cross sections. This extrapolated cross section is a factor of $1.5 - 2$ higher than those on hydrogen targets [14]. An extrapolation with $\alpha = 1$ gives better agreement with the hydrogen target data, at least for $\Xi^-$ production.

We employ the intrinsic model [23, 24, 25, 26, 27] as developed for strangeness pro-

\footnote{For the two parameterizations to be equivalent, the right-hand side of Eq. (1) should be multiplied by $2E/\sqrt{s}$ to obtain the invariant cross section.}
duction in Ref. [28]. In the intrinsic model, a hadron can fluctuate into Fock state configurations with a combination of light and strange quark pairs. The heavier quarks in the configuration are comoving with the other partons in the Fock state and thus can coalesce with these comoving partons to produce strange hadrons at large $x_F$. The model combines leading-twist production of $s\bar{s}$ pairs with intrinsic Fock states with up to nine particles. Thus coalescence production of the Ω from a proton is possible.

2 Leading-Twist Production

We calculate leading-twist strangeness in perturbative QCD, assuming the strange quark is massive. When the projectile has nonzero strangeness, we also consider the possibility of flavor excitation. We choose proton parton distribution functions with the lowest possible initial scale $\mu_0^2$ so that $m_s^2 > \mu_0^2$. Therefore the baryon parton distribution functions are based on the GRV 94 LO proton parton distributions [29]. We use the most recent pion parton densities by Glück, Reya and Schienbein [30]. To be conservative, we assume that the scale $\mu$ at which the strong coupling constant and the parton densities are evaluated is $\mu = 2m_T$ where $m_T = \sqrt{p_T^2 + m_s^2}$ and $m_s = 500$ MeV. The $x_F$ distributions, obtained by integrating Eq. (1) or (3) over $p_T$, are dominated by low $p_T$ production.

We treat the strange quark as heavy, as in Ref. [28], rather than as a massless parton in jet-like processes. Treating the strange quark as a jet means selecting a minimum $p_T$ to keep the cross section finite. A large minimum $p_T$ compatible with hard scattering is incompatible with the assumption of intrinsic production, inherently a low $p_T$ process [25]. Strange hadrons can either be produced directly in jet production or by the fragmentation of light quark and gluon jets. However, there is no indication that these data originate from jets.

The leading-twist $x_F$ distribution of heavy quark production [31] is denoted by $F$,

$$ F \equiv \frac{d\sigma_{ht}^S}{dx_F} = \frac{\sqrt{s}}{2} \int d\alpha_3 dy_2 dp_T^2 x_a x_b H_{AB}(x_a, x_b, \mu^2) \frac{1}{E_1} \frac{D_{S/s}(z_3)}{z_3}, \quad (4) $$

where $A$ and $B$ are the initial hadrons, $a$ and $b$ are the interacting partons, 1 and 2 are the produced strange quarks and 3 is the final-state strange hadron $S$. The $x_F$ of the detected quark is $x_F = 2m_T \sinh y/\sqrt{s}$ where $y$ is the rapidity of the quark and $\sqrt{s}$ is the hadron-hadron center of mass energy. We assume the simplest possible fragmentation function,

$$ D_{S/s}(z) = B_S \delta(1 - z), \quad (5) $$

with $B_S = 0.1$, assuming that all 10 ground-state strange hadrons are produced at the same rate to leading twist [28]. This choice of $D_{S/s}$ gives the hardest possible leading twist $x_F$ distribution.
The convolution of the subprocess $q\bar{q}$ annihilation and gluon fusion cross sections with the parton densities is included in $H_{AB}(x_a, x_b, \mu^2)$,

$$H_{AB}(x_a, x_b, \mu^2) = \sum_q (f_q^A(x_a, \mu^2) f_{\bar{q}}^B(x_b, \mu^2) + f_{\bar{q}}^A(x_a, \mu^2) f_q^B(x_b, \mu^2)) \frac{d\hat{\sigma}_{q\bar{q}}}{dt}$$

(6)

$$+ f_q^A(x_a, \mu^2) f_{\bar{q}}^B(x_b, \mu^2) \frac{d\hat{\sigma}_{gg}}{dt},$$

where $q = u, d$, and $s$. Although including the $s$ quark in the sum over $q$ in Eq. (6) could lead to some over counting, the strange quark contribution to $F$ from non-strange projectiles is negligible, less than 0.01% for neutron and pion projectiles. It is somewhat larger for strange projectiles, 2.5% for the $\Sigma^-$ and 5.6% for the $\Xi^-$ but it is only significant at large $x_F$. Hyperon parton distributions can be inferred from the proton distributions [27] by simple counting rules. We can relate the valence $s$ distribution of the $\Sigma^-$, $f_{s_{\Sigma^-}}^\Sigma$, to the proton valence $d$ distribution, $f_{d_p}^p$, and the valence $d$ distribution in the $\Sigma^-$, $f_{d_{\Sigma^-}}^\Sigma$, to the valence $u$ in the proton, $f_{u_p}^p$, so that

$$\int_0^1 dx f_{\Sigma^-}^\Sigma(x, \mu^2) = \int_0^1 dx f_{d_p}^p(x, \mu^2) = 1,$$

(7)

$$\int_0^1 dx f_{d_{\Sigma^-}}^\Sigma(x, \mu^2) = \int_0^1 dx f_{u_p}^p(x, \mu^2) = 2.$$

(8)

We also identify the up quark in the sea of the $\Sigma^-$ with the strange quark in the proton sea, $f_{s_{\Sigma^-}}^\Sigma(x, \mu^2) = f_s^p(x, \mu^2)$. Similar relations hold for the antiquark distributions. Likewise, for the $\Xi^-$, we relate the valence $s$, $f_{s_{\Xi^-}}^\Xi$, to the valence $u$ in the proton, $f_{u_p}^p$, and equate the valence $d$ distributions in both baryons so that,

$$\int_0^1 dx f_{s_{\Xi^-}}^\Xi(x, \mu^2) = \int_0^1 dx f_{u_p}^p(x, \mu^2) = 2,$$

(9)

$$\int_0^1 dx f_{d_{\Xi^-}}^\Xi(x, \mu^2) = \int_0^1 dx f_{d_p}^p(x, \mu^2) = 1.$$

(10)

Here also, $f_{u_p}^p(x, \mu^2) = f_p^p(x, \mu^2)$. The gluon distributions are assumed to be identical for all baryons, $f_g^p = f_g^{\Sigma^-} = f_g^{\Xi^-}$. The leading order subprocess cross sections for heavy quark production can be found in Ref. [32]. The fractional momenta carried by the projectile and target partons, $x_a$ and $x_b$, are $x_a = (m_T/\sqrt{s})(e^y + e^{\mu_2})$ and $x_b = (m_T/\sqrt{s})(e^{-y} + e^{-\mu_2})$.

We have assumed only $gg$ and $q\bar{q}$ contributions to massive strange quark production. We have also checked how the $x_F$ distribution would change if the strange quark was treated as massless and all $2 \to 2$ scattering channels were included. Jet production of $s$ quarks is through processes such as $gs \to gs$, $qs \to qs$ and $q\bar{s} \to q\bar{s}$. (Similarly for the $\Sigma$.) Including these processes increases the cross section by a factor of $4 - 8$. While this factor is not constant, it increases rather slowly with $x_F$ so that the difference in shape is only important in the region where intrinsic production dominates, as discussed later.
Contributions from massless $2 \to 2$ scattering increase more rapidly at $x_F > 0$ for strange projectiles because the contribution from, for example, $f_{s-}^\Sigma(x_a)f_p^p(x_b)$, dominates the scattering cross section. In the infinite momentum frame, $f_{s-}^\Sigma = f_{d-}^p$, see Eq. (8), and $f_{s-}^\Sigma$ is large at large $x_a$ while $f_p^p$ increases as $x_b$ decreases. To take this into account quantitatively, we have incorporated “flavor excitation” of massive strange valence quarks. The excitation matrix elements for massive quarks are found in Ref. [33].

The flavor excitation cross section has a pole when $p_T \to 0$ so that a cutoff, $p_{T\text{min}}$, is required to keep this cross section finite, as in jet production. We employ $p_{T\text{min}} = 2m_s = 1$ GeV. The leading-twist fusion cross section for strange projectiles is then augmented by

$$X_{p_{T\text{min}}} \equiv \frac{d\sigma_{\text{lt}e}}{dx_F} = \frac{\sqrt{s}}{2} \int dz_3 dy_2 dp_T^2 x_a' x_b' H'_{AB}(x_a', x_b', \mu^2) \frac{D_{S/s}(z_3)}{z_3}$$

where

$$H'_{AB}(x_a', x_b', \mu^2) = f^A_{s-}(x_a', \mu^2)\{\sum_q [f^B_q(x_b', \mu^2) + f^B_{q\bar{q}}(x_b', \mu^2)] \frac{d\hat{\sigma}_{sq}}{dt} + f^B(x_b', \mu^2) \frac{d\hat{\sigma}_{sg}}{dt}\} ,$$

$x_a' = (m_T e^{y} + p_T e^{y_2})/\sqrt{s}$ and $x_b' = (m_T e^{-y} + p_T e^{-y_2})/\sqrt{s}$. Note that there is no overlap between the processes included in Eqs. (12) and (6) and thus no double counting. This excitation mechanism is effective only for hadrons with a strange quark in the final state and thus does not affect the distributions with a produced $\bar{s}$.

To summarize, for strange and antistrange final states produced by non-strange hadrons,

$$\frac{d\sigma^S_{\text{lt}e}}{dx_F} \equiv F ,$$

as in Eq. (4). This relation also holds for antistrange final states from strange hadrons. However, for strange hadron production by hadrons with nonzero strangeness, we also consider

$$\frac{d\sigma^S_{\text{lt}e}}{dx_F} = \frac{d\sigma^S_{\text{ltf}}}{dx_F} + \frac{d\sigma^S_{\text{lt}e}}{dx_F} \equiv F + X_{p_{T\text{min}}},$$

where flavor excitation, Eq. (11), may play a role.

We remark that the role of flavor excitation in heavy quark production, as outlined in Ref. [33], is questionable. It was first proposed as a leading order contribution to the total cross section and, as such, could be rather large if the heavy quark distribution in the proton is significant. However, the proton heavy quark distribution is only nonzero above the threshold $m_Q$. In addition, parton distribution functions are defined in the infinite momentum frame where the partons are treated as massless. Later studies at next-to-leading order (NLO) [34] showed that these excitation diagrams are a subset of the NLO cross section and suppressed relative to fusion production. They are only a small fraction of the heavy flavor cross section and thus play no significant role at
Figure 1: Leading-twist strange quark production in (a) $\pi^- p$ interactions at 345 GeV, (b) $np$ interactions at 260 GeV, (c) $\Sigma^- p$ interactions at 345 GeV, and (d) $\Xi^- p$ interactions at 116 GeV. The solid curves are the fusion results, $F$. For projectiles with valence strange quarks, the excitation contributions, $X_{pT_{\text{min}}}$ with $pT_{\text{min}} = 1$ GeV, are shown in the dashed curves. The dot-dashed curves are the total, $F + X_{pT_{\text{min}}}$. 

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\end{figure}
low energies. Strange hadron production at large $x_F$ is then an important test of the excitation process.

The leading-twist $x_F$ distributions with all four projectiles are shown in Fig. 1. We compare the fusion $x_F$ distributions, $F$, in $\pi^-p$ and $\Sigma^-p$ interactions at 345 GeV, $np$ interactions at 260 GeV, and $\Xi^-p$ interactions at 116 GeV, corresponding to the energies we investigate. The $\pi^-p$ distribution is the broadest because the $f_{\pi^-u}^p(x_a)f_{u}^{p}(x_b)$ contribution hardens the $x_F$ distribution at large $x_F$ where the $q\bar{q}$ channel dominates. When excitation is considered, as in $\Sigma^-p$ and $\Xi^-p$ interactions, the $x_F$ distribution is hardened, particularly through the $f_{\Sigma^-s}^p(x'_a)f_{g}^{p}(x'_b)$ and $f_{\Xi^-s}^{p}(x'_a)f_{g}^{p}(x'_b)$ contributions. These dominate at large $x_F$ where $f_{\Sigma^-s}^p(x'_a)$ is large for valence strange quarks and $f_{g}^{p}(x'_b)$ is large at small $x'_b$. The effect is even stronger for the $\Xi^-$ since it has two valence strange quarks. Note that, except at small $x_F$, the total leading-twist cross section is equivalent to $X_{p_{T\min}}$.

We can obtain an approximate estimate of the exponent $a$ from Eq. (1) from the average $x_F$, $\langle x_F \rangle$, where

$$a = \frac{1}{\langle x_F \rangle} - 2.$$  \hspace{1cm} (15)

When averaged over $x_F > 0$, the values of $a$ obtained are larger than those measured by WA89, as expected. For the pion and neutron beams, $a = 5.2$ and 8.7 respectively. Strangeness production by strange hadrons including fusion alone also gives large values of $a$, 9 for the $\Sigma^-$ and 7.4 for the $\Xi^-$. The $x_F$ distribution of strange quarks produced by flavor excitation is considerably harder, $a = 3.2$ for the $\Sigma^-$ and 2.2 for the $\Xi^-$. Combining the two contributions, as in the dot-dashed curves in Fig. 1(c) and (d), gives a somewhat larger value of $a$ than for flavor excitation alone, $a = 4.7$ and 2.9 for $\Sigma^-$ and $\Xi^-$ beams respectively. The values of $a$ obtained from Eq. (15) are all much higher than those obtained from the data. Thus the leading twist results alone cannot explain the shape of the measured $\Xi^- x_F$ distributions.

3 The Intrinsic Model for Strangeness

We now briefly discuss the intrinsic model for strangeness production, described in detail for $\pi^-p$ interactions in Ref. [28]. Since all the data is at $x_F > 0$, we only discuss intrinsic production in the projectile.

The hadron wavefunction is a superposition of Fock state fluctuations in which the hadron contains one or more “intrinsic” $Q\bar{Q}$ pairs. These pairs can hadronize when the hadron interacts, breaking the coherence of the state. The model, first developed for charm [23, 24], gives heavy quarks a larger fraction of the projectile momentum due to their greater mass. The strange quark is lighter so that the momentum gained is
not as large. However, the intrinsic strangeness probability is larger, $P_5^{\text{is}} \sim 2\%$. For simplicity, we assume that the intrinsic probabilities are independent of the valence quark content of the projectile. Then $P_5^{\text{is}}$ is identical for nucleons and hyperons. The Fock state probabilities for up to $3Q\overline{Q}$ pairs where at least one $Q\overline{Q}$ pair is strange are given in Ref. [28].

In this paper, we focus on $\Xi^-(dss)$ and $\Omega(sss)$ production from $\pi^-(wd)$, $n(udd)$, $\Sigma^-(dds)$ and $\Xi^-(dss)$ projectiles. The produced $\Xi^-$ shares one or more valence quarks with the projectile. We study $\Omega$ production only by $\Xi^-$ projectiles, with two valence quarks in common.

Once the coherence of the Fock state is broken, the partons in the state can hadronize in two ways. The first, uncorrelated fragmentation of the strange quark, is the same basic mechanism as at leading twist. However, when the Fock state fluctuation includes all the valence quarks of the final-state hadron, these quarks, in close spatial proximity, can coalesce into the final-state hadron and come on shell. Thus, to calculate the full strange and antistrange hadron $x_F$ distributions in the intrinsic model, we include uncorrelated fragmentation of the strange quark in every state considered and coalescence from those states where it is possible. In Ref. [28], a comparison of the model with strange baryon asymmetries suggested that fragmentation may not be an effective mechanism because when the Fock state has minimal invariant mass, fragmentation may cost too much energy. This conclusion needs to be checked against inclusive $x_F$ distributions over a broader $x_F$ range.

In principle, the parton distributions of the hadron can be defined through such a Fock-state expansion [35]. In each fluctuation, only the mass distinguishes the light and heavy quark distributions. Thus it is not really possible in a given state to separate the “valence” and “sea” distributions. All are similar as long as the quarks are light. One distinguishing feature is our assumption that only strange quarks can produce strange final-state hadrons by uncorrelated fragmentation. Thus with hyperon projectiles, uncorrelated fragmentation may also be possible from Fock states with only light $Q\overline{Q}$ pairs. These intrinsic light quark states must be included in the total probability, as described in Ref. [28]. The probabilities for these states must also be defined. We assume

$$P_5^{\text{is}} = \left( \frac{\bar{m}_s}{\bar{m}_q} \right)^2 P_5^{\text{is}} \approx 5\% ,$$  \hspace{1cm} (16)$$

$$P_7^{\text{isqq}} = \left( \frac{\bar{m}_s}{\bar{m}_q} \right)^2 P_7^{\text{isqq}} = 1.75 P_5^{\text{is}} ,$$  \hspace{1cm} (17)$$

$$P_9^{\text{issq}} = \left( \frac{\bar{m}_s}{\bar{m}_q} \right)^4 P_9^{\text{issq}} = 1.25 P_5^{\text{is}} .$$  \hspace{1cm} (18)$$

We further assume that the probabilities for the meson Fock configurations are equal to the baryon probabilities.

We have only taken the 10 strange ground state hadrons and antihadrons into ac-
count. We assume that each hadron has a 10% production probability from fragmentation, neglecting the particle masses. The final-state $x_F$ distribution is then equivalent to that of the $s$ or $\bar{s}$ quark. For coalescence, we count the number of possible strange and antistrange hadron combinations that can be obtained from a given Fock state and assign each strange hadron or antihadron that fraction of the total. The possible number of strange hadrons is greater than the number of possible strange antihadrons. We clearly err in the overall normalization by simply including the ground state strange particles. However, the higher-lying resonances have the same quark content with the same fragmentation and coalescence distributions since all properties of the final-state hadrons except their quark content are neglected.

To obtain the total probability of each strange hadron to be produced from projectile hadron, $h$, in the intrinsic model, we sum the probabilities over all the states with up to 3 intrinsic $Q\bar{Q}$ pairs. Thus

$$\frac{dP^h_S}{dx_F} = \sum_n \sum_{m_u} \sum_{m_d} \sum_{m_s} \beta \left( \frac{1}{10} \frac{dP^{nF}_{i(m_s)(m_u)(m_d)}}{dx_F} + \xi \frac{dP^{nC}_{i(m_s)(m_u)(m_d)}}{dx_F} \right). \quad (19)$$

To conserve probability, $\beta = 1$ when the hadron is only produced by uncorrelated fragmentation and 0.5 when both fragmentation and coalescence are possible. When we assume coalescence production only, we set $P^{nF} \equiv 0$ and $\beta \equiv 1$. The weight of each state produced by coalescence is $\xi$ where $\xi = 0$ when $S$ is not produced by coalescence in state $|n, m_s(s\bar{s})m_u(u\bar{u})m_d(d\bar{d})\rangle$. The number of up, down and strange $Q\bar{Q}$ pairs is indicated by $m_u$, $m_d$ and $m_s$ respectively. The total, $m_u + m_d + m_s = m$, is defined as $m = (n - n_v)/2$ because each $Q$ in an $n$-parton state is accompanied by a $\bar{Q}$. For baryon projectiles, $n = 5, 7$, and 9 while for mesons $n = 4, 6$, and 8. Depending on the value of $n$, $m_i$ can be 0, 1, 2 or 3, e.g. in a $|uuds\bar{s}d\bar{s}d\bar{d}\rangle$ state, $m_u = 0$, $m_d = 2$ and $m_s = 1$ with $m = 3$. Note that $m_s = 0$ is only possible when $h$ is strange since no additional $s\bar{s}$ pairs are thus needed to produce some strange hadrons by coalescence. The total probability distributions, $dP^h_S/dx_F$, for $S = \Xi^-$ and $\Omega$ are given in the Appendix.

### 4 A Dependence of Combined Model

The total $x_F$ distribution for final-state strange hadron $S$ is the sum of the leading-twist fusion and intrinsic strangeness components,

$$\frac{d\sigma^S_{iN}}{dx_F} = \frac{d\sigma^S_U}{dx_F} + \frac{d\sigma^S_{iQ}}{dx_F}. \quad (20)$$

The leading-twist distributions are defined in Eqs. (13) and (14). The total intrinsic cross section, $d\sigma^S_{iQ}/dx_F$, is related to $dP^h_S/dx_F$ by

$$\frac{d\sigma^S_{iQ}}{dx_F} = \sigma^{in}_{hN} \frac{\mu^2}{4m_s^2} \frac{dP^h_S}{dx_F}. \quad (21)$$
We assume that the relative $A$ dependence for leading-twist and intrinsic production is the same as that for heavy quarks and quarkonia [25, 27, 31, 36]. The $A$ dependence of the two component model is

$$\sigma_{hA} = A^\beta \sigma_{lt} + A^\gamma \sigma_{iQ} \quad (22)$$

where the combination of the two terms should approximate $A^\alpha$ in Eq. (2). There are no strong nuclear effects on open charm at leading twist so that the $A$ dependence is linear at $x_F \sim 0$, $\alpha = 1$, [37], dropping to $\alpha = 0.77$ for pions and 0.71 for protons [31, 38] at higher $x_F$ where the intrinsic model begins to dominate. We assume that $\beta = 1$ and $\gamma = 0.77$ for pions and 0.71 for all baryons. Thus,

$$A^{\gamma-\beta} \approx A^{-1/3} \quad \text{as} \quad x_F \to 1 \quad (23)$$

This relative $A$ dependence, similar to that discussed earlier for light hadrons [19], is included in our calculations. The proton and neutron numbers are taken into account in the calculation of the leading-twist cross section. This isospin effect is small for fusion, $F$, which is dominated by the $gg$ channel. In perturbative QCD, $\beta = 1$ could be modified by changes in the nuclear parton distributions relative to the proton [39]. However, the scale for our perturbative calculation is too low for such models of these modifications to apply [40, 41] and are not included in our calculations.

5 Results

We begin by comparing the model to the WA89 pion data in Fig. 2. These data do not strongly distinguish between leading-twist fusion and the full model. The intrinsic results do not significantly depend on uncorrelated fragmentation. All three curves agree rather well with the data, primarily because the fusion $x_F$ distribution is already fairly hard. Then the intrinsic contribution is a small effect even though the $d$ valence quark is common between the $\pi^-$ and the $\Xi^-$. This is perhaps due to the fact that $P_{iss}^6$ is already rather small, $\sim 0.6\%$. We note that the calculated total cross sections agree with the measured cross sections to within better than a factor of two despite the rather large uncertainties in the calculations.

Even though the intrinsic contribution is relatively small, it significantly affects the value of $a$ obtained from Eq. (15). The difference between the $a$ values found without and with uncorrelated fragmentation in the intrinsic model is negligible for the pion beam. We find $a \approx 4.1$ for the C target and 4.3 for the Cu target relative to $a = 5.2$ for leading twist alone. These results are within the errors of the WA89 fit to their data. The agreement is especially good since the two-component model does not give a smooth falloff as a function of $x_F$ that can be easily quantified by a single exponent.

It is also possible to calculate $\alpha(x_F)$ and the $x_F$-integrated $\alpha$ from the ratios of the distributions. The calculations including both uncorrelated fragmentation and coalescence generally give a smaller value of $\alpha$ and, hence, a stronger $A$ dependence. This
Figure 2: The model is compared to the 345 GeV WA89 pion data on (a) C and (b) Cu targets. The dotted curves are leading-twist fusion, $F$, alone, the dashed curves include uncorrelated fragmentation and coalescence, and the solid curves include coalescence alone. The data sets have been normalized to the cross section per nucleon. The curves are normalized to the data at $x_F = 0.15$.

is because fragmentation peaks at low $x_F$, influencing the $A$ dependence sooner than coalescence alone which is only significant at intermediate $x_F$. Thus $\alpha(x_F) \sim 0.9$ for fragmentation and coalescence while $\alpha$ decreases from $\sim 1$ at low $x_F$ to 0.86 at high $x_F$ with coalescence alone. The integrated values are 0.93 and 0.98 respectively, somewhat higher than the WA89 result but with the same general trend.

The overall agreement with the total cross section is not as good for the $nA$ data, shown in Fig. 3. Surprisingly, the distribution including uncorrelated fragmentation agrees best with the data. This is perhaps because the energy of the secondary neutron beam is least well determined. The energy spread is 15% compared to 9% for the pion and $\Sigma^{-}$ beams. A small energy variation can have a large effect on the leading-twist cross section. A 15% increase in the average neutron energy, from 260 GeV to 300 GeV, increases $d\sigma_{lt}/dx_F$ by 80% at $x_F \sim 0.25$ while the intrinsic cross section is essentially unaffected. Such a shift in the relative leading-twist and intrinsic production rates could easily reduce the effect of coalescence alone to be more compatible with the data. The uncertainty in the energy of the pion beam has a much weaker effect on the final result because the intrinsic contribution is already small, as is obvious from Fig. 2.

The calculated exponents $a$ are larger for the neutron than the pion, in agreement with the WA89 measurements [17]. We find $a \approx 4.9$ for C and 5.8 for Cu. Typically the value of $a$ obtained for fragmentation and coalescence is larger than that for coalescence alone since eliminating the fragmentation contribution tends to increase $\langle x_F \rangle$. The stronger $A$ dependence assumed for the intrinsic model has the effect of increasing $a$ for larger nuclei. Thus the $a$ found for the carbon target agrees rather well with the WA89
Figure 3: The model is compared to the 260 GeV WA89 neutron data on (a) C and (b) Cu targets. The dotted curves are leading-twist fusion, $F$, alone, the dashed curves include uncorrelated fragmentation and coalescence, and the solid curves include coalescence alone. The data sets have been normalized to the cross section per nucleon. The curves are normalized to the data at $x_F = 0.15$.

data while the copper data suggest a harder distribution than our calculation implies. There is, however, a stronger $A$ dependence in the falloff with $x_F$ than in the data which cannot distinguish between the values of $a$ determined for the two targets. This stronger dependence is reflected in the values of $\alpha$ obtained, 0.87 when fragmentation and coalescence are included and 0.95 with coalescence alone.

Even though there is some $A$ dependence in the model calculations, the relatively small intrinsic contribution to the pion and neutron data leads to a rather weak overall $A$ dependence. Dominance of the leading-twist cross section at low to intermediate $x_F$ results in a nearer-to-linear integrated $\alpha$, as observed by WA89 [17].

We now turn to $\Xi^-$ production by the $\Sigma^-$ where the $A$ dependence can be expected to be stronger. For the first time, we have a valence strange quark in the projectile so that we can compare the effectiveness of fusion alone with flavor excitation. We can also better test the importance of uncorrelated fragmentation because coalescence production is already possible in the 5-parton Fock state, $|ddss\bar{s}\rangle$.

Our results are collected in Fig. 4. We first discuss the importance of uncorrelated fragmentation to leading-twist fusion, $F$, alone, Fig. 4(a) and (b). The leading-twist contribution is rather steeply falling. Including both uncorrelated fragmentation and coalescence broadens the $x_F$ distribution but cannot match the hardness of the measured $x_F$ distribution. Eliminating the fragmentation contribution produces a much harder distribution for $x_F \geq 0.15$, matching the shape of the data relatively well.
Figure 4: The model is compared to the 345 GeV WA89 $\Sigma^- A$ data on C and Cu targets. In (a) and (b), the leading twist contribution is $F$ while in (c) and (d), flavor excitation is also included, $F + X_{pT_{min}}$. The dotted curves are for leading-twist alone, the dashed curves include uncorrelated fragmentation and coalescence and the solid curves include coalescence alone. The data sets have been normalized to the cross section per nucleon. The curves are normalized to the data at $x_F = 0.15$. 

$\sigma / dx_F \, (\mu b)$

$X_F$

(a) C

(b) Cu

(c) C

(d) Cu
The calculations are all normalized to the $x_F = 0.15$ point to better compare the shapes of the distributions. Fragmentation gives better agreement at low $x_F$ because uncorrelated fragmentation peaks close to $x_F \sim 0$, filling in the low to intermediate $x_F$ range. Coalescence, on the other hand, always produces strange hadrons with $\langle x_F \rangle \geq 0.3$, broadening the distribution only in this region. Thus without fragmentation the calculations overestimate the data at $x_F \sim 0$. The data seem to indicate that uncorrelated fragmentation is not an effective mechanism for intrinsic production, in agreement with the conclusions of Ref. [28].

The agreement with the solid curves in Fig. 4(a) and (b) is good but not perfect. The calculation overestimates the data at high $x_F$. Recall that for the neutron, the 15% spread in the beam momentum could result in an overestimate of the intrinsic contribution, as previously discussed. Although the possible spread in the $\Sigma^-$ beam momentum is smaller, it could affect the relative intrinsic contribution at low to intermediate $x_F$. At large $x_F$, the effect on the shape would be negligible because the intrinsic contribution dominates. Thus, given the inherent uncertainties in the model and in the data, the agreement is rather satisfactory.

We have obtained the value of the exponent $a$ from $\langle x_F \rangle$, both over all $x_F$ and for $x_F > 0.1$, avoiding the strong change in slope of the solid curves when coalescence alone is included in the intrinsic result. When the entire forward $x_F$ range is integrated over, $a = 3.02$ for C and 3.39 for Cu with both uncorrelated fragmentation and coalescence while with coalescence alone, $a = 1.24$ for C and 1.78 for Cu. Considering only the range $x_F > 0.1$, we find $a = 1.56$ for C and 1.63 for Cu with fragmentation and 0.43 for C, 0.57 for Cu without fragmentation. The calculated $a$'s suggest considerably harder $x_F$ distributions in the more limited $x_F$ region, particularly when coalescence alone is considered. However, none of the results are in good agreement with $a \approx 2$, as obtained by WA89 for $x_F \geq 0.4$. This is not surprising, especially since the solid curve is seen to be harder than the data for $x_F > 0.1$. In any case, the coalescence contributions in particular, now considerably more important than for $\Xi^-$ production by non-strange hadrons, are difficult to fit to a power law since they approach zero at both $x_F = 0$ and $x_F = 1$ with a peak at intermediate $x_F$, see the curves in Ref. [28]. The various contributions, all with somewhat different magnitudes due to the relative probabilities, peak at different values of $x_F$, complicating the situation further.

The calculated values of $\alpha$ give $\alpha \approx 0.8$ for the integrated cross sections but $\alpha \approx 0.7$ for $x_F > 0.4$, with and without fragmentation in the intrinsic model, rather consistent with the WA89 result. However, as a function of $x_F$, coalescence alone is more consistent with the measurements since $\alpha \approx 1$ at $x_F \approx 0$, decreasing to 0.71 as $x_F \rightarrow 1$, as expected from Eq. (22).

We now check if our results improve when we include flavor excitation, Eq. (14), shown in Fig. 4(c) and (d). Now the baseline leading twist distribution, $F + X_{p_{T\text{min}}}$, is harder than with fusion alone, as shown in Fig. 1(c). However, although the distribution is broader, it still drops six orders of magnitude over the entire $x_F$ range with $p_{T\text{min}} = 1$
GeV. Thus including flavor excitation cannot describe the data without the intrinsic coalescence component, as in Fig. 4(a) and (b). The total cross section is in reasonable agreement with that measured by WA89. Decreasing $p_{T\text{min}}$ further can harden the distribution but still underestimates the data. A lower $p_{T\text{min}}$ enhances the total cross sections considerably so that, if $p_{T\text{min}} = 0.25$ GeV, the cross section is overestimated by several orders of magnitude. The intrinsic contribution is then negligible so that decreasing $p_{T\text{min}}$ actually degrades the agreement with the data. Thus there is no clear evidence for flavor excitation.

The trends in the $A$ dependence are similar when excitation is included although the values of $a$ obtained are suggestive of harder $x_F$ distributions than with leading-twist fusion alone. In particular, the excitation contribution is harder at low $x_F$, see Fig. 1, causing the change in slope due to the hardening of the intrinsic distribution when coalescence alone is included to be less abrupt. Nonetheless, the agreement with the measured value of $a$ is not significantly improved. The calculated values of $\alpha(x_F)$ are similar to those with leading-twist fusion alone but the integrated values of $\alpha$ are somewhat larger, $\approx 0.87$, due to the larger leading-twist baseline. The $A$ dependence also does not support flavor excitation as a significant contribution to strange hadron production.

To summarize, the $A$ dependence of $\Xi^-$ production by $\Sigma^-$ is stronger because the intrinsic contribution with coalescence dominates the $x_F$ distribution already at $x_F \sim 0.1$. Therefore the integrated $A$ dependence is nearly a factor of $A^{1/3}$ down relative to the pion and neutron data, as shown in Eqs. (22) and (23). Thus the trends of the model are in agreement with the WA89 data.

Finally, we compare our intrinsic model calculations with the invariant $\Xi^-$ and $\Omega$ cross sections measured in $\Xi^-$Be interactions at 116 GeV [14]. Since intrinsic production is expected to be a primarily low $p_T$ effect [25], we only compare to the low $p_T^2$ bin, $0 < p_T^2 < 0.4$ GeV$^2$ [42]. The data and our calculations are shown in Fig. 5. We have multiplied our $x_F$ distributions by $2m_T \cosh y/\sqrt{s}$ to obtain the invariant cross section. The invariant $x_F$ distributions are harder as a function of $x_F$.

Because the initial and final states are identical for $\Xi^-$ production, the intrinsic contribution increases with $x_F$. However, even with coalescence alone, the increase does not continue beyond $x_F \sim 0.4$. A similar but weaker effect is seen for the $\Omega$ where there are two strange quarks in common with the $\Xi^-$. Therefore we have tried to identify a mechanism that would increase the cross section beyond $x_F \sim 0.4$. One possibility is a “Pomeron-like” parton in the Fock state. Since the Pomeron has quantum numbers similar to two gluons, it can be exchanged between two projectile valence quarks. A $|dssP\rangle$ state, where $P$ signifies the “Pomeron”, would result in $\Xi^-$ states at high $x_F$ while avoiding the $\delta(1 - x_F)$ delta function for the 3-particle Fock state. A $\Xi^-$ from such a configuration would have a distribution peaking at $x_F \to 1$. It is also possible to imagine a $|dss\Xi P\rangle$ state from which both the $\Xi^-$ and $\Omega$ could be produced. In this case, the distribution would peak away from $x_F \sim 1$. We included “Pomeron” production
Figure 5: The model is compared to the 116 GeV Ξ−Be data. In (a) and (b), the leading twist contribution is $F$ while in (c) and (d), flavor excitation is also included, $F + X_{p_{T_{\text{min}}}}$. The dotted curves are for leading-twist fusion alone, the dashed curves include uncorrelated fragmentation and coalescence, the solid curves include coalescence alone and the dot-dashed curves include a diffractive “Pomeron” contribution. The data sets have been normalized to the cross section per nucleon. The curves are normalized to the data at $x_F = 0.15$. 
from both states assuming that $P_{4}^{i}P_{4} = P_{5}^{i} \sim 5\%$ and that $P_{6}^{i}P_{6} = P_{7}^{i} \sim 3.5\%$, giving these configurations large probability. The results, shown in the dot-dashed curves in Fig. 5, agree relatively well with the data, especially for the $\Xi^{-}$. The $\Omega$ data are still underestimated but the trend is now in the right direction.

Of course, this is a rather artificial solution, especially when the initial and final states are not identical. If it is correct, it should also be included in $\Sigma^{-}p \rightarrow \Xi^{-}X$, as shown in Fig. 4. However, we have checked this and found that the resulting $x_F$ distribution is far too hard. Therefore, the practicality of the mechanism is questionable and the “Pomeron” results should not be taken too seriously.

Our calculations with flavor excitation are compared to the data in Fig. 5(c) and (d). The results do not improve, even when the “Pomeron” is included. Indeed, the results with excitation underestimate the data at high $x_F$ more than with fusion alone. Therefore we conclude that flavor excitation is not an effective mechanism for strange hadron production at low $p_T$, in keeping with the interpretation of the excitation diagrams as NLO contributions to the production cross section, as discussed previously.

We have also calculated the exponent $a'$, see Eq. (3), for these distributions with $x_F > 0.5$. The results are negative for all the cases shown. We find $a' \approx -0.45$ with and without uncorrelated fragmentation and $-0.5$ with the “Pomeron”. These values are rather consistent with those obtained for the low $p_T^2$ selection of the $\Xi^{-}$ data. The corresponding values for $\Omega$ production are somewhat lower, $\approx -0.41$, without the “Pomeron” but somewhat higher, $\approx -0.54$, with it.

6 Conclusions

We have compared our intrinsic calculations to $\Xi^{-}$ production by $\pi^{-}$, $n$ and $\Sigma^{-}$ projectiles and to $\Xi^{-}$ and $\Omega$ production by $\Xi^{-}$ projectiles. We find good agreement with the WA89 data for leading-twist fusion and coalescence. Flavor excitation seems excluded as a significant mechanism of low $p_T$ strange hadron production. The apparent difficulties with uncorrelated fragmentation seen in Ref. [28] are confirmed here. Thus coalescence production is the most effective mechanism for strange hadron production in the intrinsic model. The leading charm analysis should perhaps be revisited in light of this conclusion.

The conclusions that can be reached from the $\Xi^{-}$-induced interactions at 116 GeV are less clear. It is possible that a “Pomeron-like” state could exist in the hadron wavefunction but its applicability to $\Omega$ production is somewhat doubtful. Therefore the interpretation of these data within the intrinsic model is rather inconclusive. More standard studies of diffractive production in $\Xi^{-}$Be $\rightarrow \Xi^{-}X$ should be performed.
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Appendix

Here we give the relevant probability distributions in the intrinsic model for $\Xi^-$ and $\Omega$ production used in our calculations. To more clearly distinguish between the probability distributions including uncorrelated fragmentation and coalescence and those with coalescence alone, both distributions are given.

First, we give the distributions relevant to the WA89 measurements. We reproduce the $\Xi^-$ probability distribution from a $\pi^-$ [28],

$$
\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{10} \frac{dP^{4F}_{\text{is}}}{dx_F} + \frac{1}{10} \frac{dP^{6F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{6F}_{\text{isd}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{6F}_{\text{iss}}}{dx_F} + \frac{1}{7} \frac{dP^{6C}_{\text{is}}}{dx_F} \right) \\
+ \frac{1}{10} \frac{dP^{8F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{8F}_{\text{isd}}}{dx_F} + \frac{1}{10} \frac{dP^{8F}_{\text{issd}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{8F}_{\text{iss}}}{dx_F} + \frac{1}{12} \frac{dP^{8C}_{\text{is}}}{dx_F} \right) \\
+ \frac{1}{2} \left( \frac{1}{10} \frac{dP^{8F}_{\text{issd}}}{dx_F} + \frac{2}{12} \frac{dP^{8C}_{\text{issd}}}{dx_F} \right) + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{8F}_{\text{iss}}}{dx_F} + \frac{3}{16} \frac{dP^{8C}_{\text{iss}}}{dx_F} \right),
$$

(1)

$$
\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{7} \frac{dP^{6C}_{\text{is}}}{dx_F} + \frac{1}{12} \frac{dP^{8C}_{\text{isu}}}{dx_F} + \frac{2}{12} \frac{dP^{8C}_{\text{isd}}}{dx_F} + \frac{3}{16} \frac{dP^{8C}_{\text{iss}}}{dx_F} .
$$

(2)

From a neutron projectile,

$$
\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{10} \frac{dP^{5F}_{\text{is}}}{dx_F} + \frac{1}{10} \frac{dP^{7F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{7F}_{\text{isd}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{7F}_{\text{iss}}}{dx_F} + \frac{1}{13} \frac{dP^{7C}_{\text{is}}}{dx_F} \right) \\
+ \frac{1}{10} \frac{dP^{9F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{9F}_{\text{isd}}}{dx_F} + \frac{1}{10} \frac{dP^{9F}_{\text{issd}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{9F}_{\text{iss}}}{dx_F} + \frac{2}{22} \frac{dP^{9C}_{\text{is}}}{dx_F} \right) \\
+ \frac{1}{2} \left( \frac{1}{10} \frac{dP^{9F}_{\text{issd}}}{dx_F} + \frac{3}{22} \frac{dP^{9C}_{\text{issd}}}{dx_F} \right) + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{9F}_{\text{iss}}}{dx_F} + \frac{6}{28} \frac{dP^{9C}_{\text{iss}}}{dx_F} \right),
$$

(3)

$$
\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{13} \frac{dP^{7C}_{\text{is}}}{dx_F} + \frac{2}{22} \frac{dP^{9C}_{\text{isu}}}{dx_F} + \frac{3}{22} \frac{dP^{9C}_{\text{isd}}}{dx_F} + \frac{6}{28} \frac{dP^{9C}_{\text{iss}}}{dx_F} .
$$

(4)

The $\Xi^-$ distribution from a $\Sigma^-$ projectile is,

$$
\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{10} \frac{dP^{5F}_{\text{is}}}{dx_F} + \frac{1}{10} \frac{dP^{5F}_{\text{isu}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{5F}_{\text{iss}}}{dx_F} + \frac{2}{6} \frac{dP^{5C}_{\text{is}}}{dx_F} \right) + \frac{1}{10} \frac{dP^{7F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{7F}_{\text{isd}}}{dx_F} \\
+ \frac{1}{10} \frac{dP^{9F}_{\text{isu}}}{dx_F} + \frac{1}{10} \frac{dP^{9F}_{\text{isd}}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{9F}_{\text{iss}}}{dx_F} + \frac{1}{13} \frac{dP^{9C}_{\text{is}}}{dx_F} \right) + \frac{1}{2} \left( \frac{1}{10} \frac{dP^{9F}_{\text{issd}}}{dx_F} + \frac{3}{13} \frac{dP^{9C}_{\text{issd}}}{dx_F} \right).
$$
We now present the relevant probability distributions for $\Xi^-$ and $\Omega$ production from a $\Xi^-$ projectile. First we give the $\Xi^-$ distributions,

$$\frac{dP_{\Xi^-}}{dx_F} = \frac{1}{2} \left( \frac{1}{10 \ dx_F} + \frac{1}{6 \ dx_F} \right) + \frac{1}{2} \left( \frac{1}{10 \ dx_F} + \frac{2}{6 \ dx_F} \right),$$

$$\frac{dP_{\Xi^-}}{dx_F} = \frac{2}{22 \ dx_F} + \frac{1}{22 \ dx_F} + \frac{1}{28 \ dx_F} + \frac{1}{32 \ dx_F}.$$
Finally, we give the $\Omega$ distribution from a $\Xi^-$ projectile,

\[
\frac{dP_{\Omega}^{\Xi^-}}{dx_F} = \frac{1}{10} \frac{dP_{5s_i}^{5f}}{dx_F} + \frac{1}{10} \frac{dP_{5s_i}^{5f}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP_{i}^{5f}}{dx_F} + \frac{1}{7} \frac{dP_{i}^{5f}}{dx_F} \right) + \frac{1}{10} \frac{dP_{i}^{5f}}{dx_F} + \frac{1}{10} \frac{dP_{i}^{5f}}{dx_F} + \frac{1}{2} \left( \frac{1}{10} \frac{dP_{i}^{5f}}{dx_F} + \frac{1}{16} \frac{dP_{i}^{5f}}{dx_F} \right),
\]

\[
\frac{dP_{\Omega}^{\Xi^-}}{dx_F} = \frac{1}{7} \frac{dP_{i}^{5c}}{dx_F} + \frac{1}{16} \frac{dP_{i}^{7c}}{dx_F} + \frac{1}{16} \frac{dP_{i}^{7c}}{dx_F} + \frac{3}{17} \frac{dP_{i}^{7c}}{dx_F} + \frac{1}{10} \frac{dP_{i}^{9c}}{dx_F} + \frac{1}{28} \frac{dP_{i}^{9c}}{dx_F} + \frac{1}{28} \frac{dP_{i}^{9c}}{dx_F} + \frac{1}{10} \frac{dP_{i}^{9c}}{dx_F} + \frac{1}{37} \frac{dP_{i}^{9c}}{dx_F}.
\]

References

[1] E.M. Aitala et al. (E791 Collab.), Phys. Lett. B371 (1996) 157.
[2] M. Aguilar-Benitez et al. (LEBC-EHS Collab.), Phys. Lett. 161B (1985) 400.
[3] M. Aguilar-Benitez et al. (LEBC-EHS Collab.), Z. Phys. C31 (1986) 491.
[4] S. Barlag et al. (ACCMOR Collab.), Z. Phys. C49 (1991) 555.
[5] M.I. Adamovich et al. (WA82 Collab.), Phys. Lett. B305 (1993) 402.
[6] G.A. Alves et al. (E769 Collab.), Phys. Rev. Lett. 72 (1994) 812.
[7] R. Werding (WA89 Collab.), in Proceedings of ICHEP94, 27th International Conference on High Energy Physics, Glasgow, Scotland (1994).
[8] M.I. Adamovich et al. (WA89 Collab.), Eur. Phys. J. C8 (1999) 593.
[9] E. Ramberg (SELEX Collab.), in Hyperons, Charm and Beauty Hadrons, Proceedings of the 2nd International Conference on Hyperons, Charm and Beauty Hadrons, Montreal, Canada, 1996, edited by C.S. Kalman et al., Nucl. Phys. B (Proc. Suppl.) 55A (1997) 173.

[10] J. Engelfried et al., in Proceedings of the 5th Workshop on Heavy Quarks at Fixed Target, Rio de Janeiro, Brazil, 2000, hep-ex/0012004.

[11] M. Bourquin et al., Nucl. Phys. B153 (1979) 13.

[12] M. Bourquin et al., Z. Phys. C5 (1980) 275.

[13] M. Bourquin and J.P. Repellin, Phys. Rep. 114 (1984) 99.

[14] S.F. Biagi et al., Z. Phys. C34 (1987) 187.

[15] O. Schneider et al., Z. Phys. C46 (1990) 341.

[16] A. Beretvas et al., Phys. Rev. D34 (1986) 53.

[17] M.I. Adamovich et al. (WA89 Collab.), Z. Phys. C76 (1997) 35.

[18] J.F. Gunion, Phys.Lett. B88 (1979) 150.

[19] W.M. Geist, Nucl. Phys. A525 (1991) 149c.

[20] T.R. Cardello et al., Phys. Rev. D32 (1985) 1.

[21] D.S. Barton et al., Phys. Rev. D27 (1983) 2580.

[22] P. Skubic et al., Phys. Rev. D18 (1978) 3115.

[23] S.J. Brodsky, P. Hoyer, C. Peterson and N. Sakai, Phys. Lett. B93 (1980) 451.

[24] S.J. Brodsky, C. Peterson and N. Sakai, Phys. Rev. D23 (1981) 2745.

[25] R. Vogt and S.J. Brodsky, Nucl. Phys. B438 (1995) 261.

[26] R. Vogt and S.J. Brodsky, Nucl. Phys. B478 (1996) 311.

[27] T. Gutierrez and R. Vogt, Nucl. Phys. B539 (1999) 189.

[28] T.D. Gutierrez and R. Vogt, Nucl. Phys. A705 (2002) 396.

[29] M. Glück, E. Reya, A. Vogt, Z. Phys. C67 (1995) 433.

[30] M. Glück, E. Reya and I. Schienbein, Eur. Phys. J. C10, 313 (1999).

[31] R. Vogt, S.J. Brodsky and P. Hoyer, Nucl. Phys. B383 (1992) 643.

[32] R.K. Ellis, in Physics at the 100 GeV Scale, Proceedings of the 17th SLAC Summer Institute, Stanford, California, 1989, edited by E.C. Brennan (SLAC Report No. 361) 45.
[33] B.L. Combridge, Nucl. Phys. B151 (1979) 429.

[34] P. Nason, S. Dawson, and R.K. Ellis, Nucl. Phys. B303 (1988) 607.

[35] P. Hoyer and D.P. Roy, Phys. Lett. B410 (1997) 63.

[36] R. Vogt, S.J. Brodsky, and P. Hoyer, Nucl. Phys. B360 (1991) 67.

[37] M.J. Leitch et al. (E789 Collab.), Phys. Rev. Lett. 72 (1994) 2542.

[38] J. Badier et al. (NA3 Collab.), Z. Phys. C20 (1983) 101.

[39] M. Arneodo, Phys. Rep. 240 (1994) 301.

[40] K.J. Eskola, V.J. Kolhinen and P.V. Ruuskanen, Nucl. Phys. B535 (1998) 351.

[41] K.J. Eskola, V.J. Kolhinen and C.A. Salgado, Eur. Phys. J. C9 (1999) 61.

[42] http://cpt19.dur.ac.uk/cgi-hepdata/hepreac/1637169