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Hyperon polarization in different inclusive production processes in unpolarized high energy hadron-hadron collisions

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We apply the picture proposed in a previous Letter, which relates the hyperon polarization in unpolarized hadron-hadron collisions to the left-right asymmetry in singly polarized reactions, to the production of different hyperons in reactions using different projectiles and/or targets. We discuss the different ingredients of the proposed picture in detail and present the results for hyperon polarization in reactions such as *pp*, K^-p, $\pi^{\pm}p$, and Σ^-p collisions. We compare the results with the available data and make predictions for future experiments.

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I. INTRODUCTION

Since its discovery in the 1970s [1,2], the surprisingly large transverse polarization of hyperons in unpolarized high energy hadron-hadron and hadron-nucleus collisions has been a standing hot topic in high energy spin physics (see, e.g., [1–13] and the references cited therein). Experimentally, there are a large number of similar experiments that have been performed at different energies and/or using different projectiles and/or targets and for the production of different hyperons [3]. Theoretically, different models have been proposed [4–13], the aim of which is to understand the origin(s) of this striking spin effect in high energy reactions. Clearly, such studies should provide us with useful information on the spin structure of the hadron and the spin dependence of strong interactions.

Inspired by the similarities of the corresponding data [3,14,15], we proposed a new approach in a recent Letter [10] to understanding the origin(s) of the transverse hyperon polarization in unpolarized hadron-hadron collisions by relating them to the left-right asymmetries observed [14,15] in singly polarized *pp* collisions. We pointed out that these two striking spin phenomena should be closely related to each other and have the same origin(s). We showed that, using the spin correlation deduced from the single-spin left-right asymmetries for inclusive π production as input, we can naturally understand the transverse polarization for a hyperon which has one valence quark in common with the projectile, such as Σ^-, Ξ^0, or Ξ^- in *pp* collisions, or Λ in K^-p collisions. We showed also that, to understand the puzzling transverse polarization of Λ in *pp* collisions, which has two valence quarks in common with the projectile, we need to assume that the *s* and $\bar{s}$, which combine, respectively, with the valence (*ud*) diquark and the remaining *u* valence quark to form the produced Λ and the associatively produced K^+ in the fragmentation region, have opposite spins. Under this assumption, we obtained a good quantitative fit to the x_F dependence of Λ polarization in *pp* collisions (where $x_F \equiv 2p_{\parallel}/\sqrt{s}$, $p_{\parallel}$ is the longitudinal component of the momentum of the produced hyperon, and $\sqrt{s}$ is the total center of mass energy of the *pp* system). The qualitative features obtained for the polarizations of other hyperons are all in good agreement with the available data.

There are two main points in the picture that need to be further tested, i.e., that (i) the *s* and $\bar{s}$ that combine, respectively, with the valence (*ud*) diquark and the remaining *u* valence quark of the projectile proton to form the produced Λ and the associatively produced K^+ in the fragmentation region have opposite spins, and (ii) the SU(6) wave function can be used to describe the relation between the spin of the fragmenting quark and that of the hadron produced in the fragmentation process. There have been developments since the publication of Ref. [10]. We found that exclusive reactions such as $pp \to p\Lambda K^+$ and $e^-p \to e^-\Lambda K^+$ are most suitable to test point (i). We therefore applied the picture to these processes and presented the results obtained in Refs. [11] and [16]. It is encouraging to see that these results are all in agreement with the available data [17,18].

It has also been pointed out [19] that the longitudinal polarization of Λ in e^+e^- annihilation at the Z^0 pole provides a special test for point (ii), i.e., whether the SU(6) wave function can be used in relating the spin of the fragmenting quark to that of the produced hadron. Calculations have been made [19,20] and the results obtained are consistent with the data [21,22]. Since neither the accuracy nor the abundance of the data is high enough to give a conclusive judgment, we made a systematic study of hyperon polarization in other reactions which can be used to test this point [19,23]. The results are in agreement with the available data [21,22,24], and future experiments are under way.

Encouraged by these developments, in this paper, we apply the proposed picture to study the polarizations of different hyperons in different unpolarized hadron-hadron and hadron-nucleus reactions. We summarize the different ingredients of the picture and their development in detail in Sec. II. In Sec. III, we give the method of calculation of hyperon polarization in different processes using the picture. In Sec. IV, we apply the method to calculate the polarizations of different hyperons produced in unpolarized *pp*, K^-p, $\pi^{\pm}p$, and Σ^-p collisions. We present the results obtained and

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compare them with the available data. Finally, a short summary and outlook are given in Sec. V.

II. THE PHYSICAL PICTURE

In this section, we summarize the key points of the picture proposed in Ref. [10]. The basic idea of the picture is that there should be a close relation between hyperon polarization (P_H) in unpolarized hadron-hadron collisions and left-right asymmetry (A_N) in single-spin hadron-hadron collisions. Hence, if we extract the essential information encoded in the A_N data, we can study P_H based on such information. There are three key points in this physical picture which are summarized in the following.

A. Correlation between the spin of the quark and the direction of motion of the produced hadron

It has been pointed out [10] that the existence of A_N in singly polarized hadron-hadron collisions implies the existence of a spin correlation between the spin of the fragmenting quark and the direction of momentum of the produced hadron, i.e., an $\vec{s}_q\cdot\vec{n}$ type of spin correlation in the reaction. [Here, $\vec{s}_q$ is the spin of the quark, $\vec{n}\equiv(\vec{p}_{inc}\times\vec{p}_h)/|\vec{p}_{inc}\times\vec{p}_h|$ is the unit vector in the direction normal to the production plane, and $\vec{p}_{inc}$ and $\vec{p}_h$ are, respectively, the momentum of the incident hadron and that of the produced hadron.] One of the major ingredients of the picture proposed in Ref. [10] is that both the existence of A_N and that of P_H are different manifestations of this spin correlation $\vec{s}_q\cdot\vec{n}$. Hence we can use the experimental results for A_N as input to determine the strength of this spin correlation, and then apply it to unpolarized hadron-hadron collisions to study P_H.

We recall that [14,15,25], in the language commonly used in describing A_N, the polarization direction of the incident proton is called upward, and the incident direction is forward. The single-spin left-right asymmetry A_N is just the difference between the cross section where $\vec{p}_h$ points to the left and that to the right, which correspond to $\vec{s}_q\cdot\vec{n}=1/2$ and $\vec{s}_q\cdot\vec{n}=-1/2$, respectively. The data on A_N show that if a hadron is produced by an upward polarized valence quark of the projectile, it has a large probability of having a transverse momentum pointing to the left [14,15]. A_N measures the excess of hadrons produced to the left over those produced to the right. The difference between the probability for the hadron to go left and that to go right is denoted by C if the hadron is produced by an upward polarized quark [25,26]. C is a constant in the range of $0<C<1$. It has been shown that [25,26], to fit the A_N data [15] in the transverse momentum interval $0.7<p_T<2.0$ GeV/c, C should be taken as $C=0.6$.

In terms of the spin correlation discussed above, the cross section should be expressed as

$$\sigma=\sigma_0+(\vec{s}_q\cdot\vec{n})\sigma_1, \tag{1}$$

where σ_0 and σ_1 are independent of $\vec{s}_q$. The second term just denotes the existence of the $\vec{s}_q\cdot\vec{n}$ type of spin correlation. C is just the difference between the cross section where $\vec{s}_q\cdot\vec{n}=1/2$ and that where $\vec{s}_q\cdot\vec{n}=-1/2$ divided by their sum, i.e., $C=\sigma_1/(2\sigma_0)$.

Now, we assume the same strength for the spin correlation in hyperon production in the same collisions. It follows that the quark which fragments into the hyperon should be polarized, and the polarization P_q can be determined by using Eq. (1). Since both the left-right asymmetry in singly polarized collisions and the hyperon polarization in unpolarized collisions exist mainly in the large x_F region, we assume that the spin correlation exists only for valence quarks of the incident hadrons. For a hyperon produced with momentum $\vec{p}_h$, $\vec{n}$ is given. The cross section for production of this hyperon in the fragmentation of a valence quark with spin satisfying $\vec{s}_q\cdot\vec{n}=1/2$ is $(\sigma_0+\sigma_1/2)$, and that for a quark with spin satisfying $\vec{s}_q\cdot\vec{n}=-1/2$ is $(\sigma_0-\sigma_1/2)$. Hence, the polarization of the valence quarks which lead to the production of hyperons with that $\vec{n}$ is given by

$$P_q=\frac{(\sigma_0+\sigma_1/2)-(\sigma_0-\sigma_1/2)}{(\sigma_0+\sigma_1/2)+(\sigma_0-\sigma_1/2)}=\frac{\sigma_1}{2\sigma_0}=C. \tag{2}$$

It should be emphasized that $P_q\neq 0$ just means that the strength of the spin correlation of the form $\vec{s}_q\cdot\vec{n}$ is nonzero in the reaction. It means that, due to some spin-dependent interactions, the quarks that have spins along the same direction as the normal to the production plane have a large probability of combining with suitable sea quarks to form the specified hyperons than those which have spins in the opposite direction. It does not imply that the quarks in the unpolarized incident hadrons were polarized in a given direction, which would contradict the general requirement of space rotation invariance. In fact, in an unpolarized reaction, the normal to the production plane of the specified hyperons is uniformly distributed in the transverse directions. Hence, averaging over all the normal directions, the quarks are unpolarized.

We would like to mention also that a similar idea has been applied to spin alignments of vector mesons in unpolarized hadron-hadron collisions [27]. It has been shown that the existence of spin alignment of the vector mesons in unpolarized hadron-hadron collision is another manifestation of the existence of the $\vec{s}_q\cdot\vec{n}$ type of spin correlation. The results obtained are in agreement with the available data [28–30].

B. Relating the spin of the quark to that of the hadron

As discussed in the above-mentioned subsection, the existence of the $\vec{s}_q\cdot\vec{n}$ type of spin correlation in hadron-hadron collisions implies a polarization of the quark q^0 transverse to the production plane of the hyperon. (Here, we use the superscript 0 to denote the quark before fragmentation.) To study the polarization of the produced hyperon from this point, we need to know the relation between the spin of the quark and that of the hadron produced in the fragmentation of this quark. The question of the relation between the spin of the fragmenting quark q^0 and that of the hadron created in the fragmentation of q^0 is usually referred to as "spin transfer in a high energy fragmentation process.'' It contains two parts: will q^0 keep its polarization in the fragmentation? and what is the relation between the spin of q^0 and that of the hadron that contains q^0?

The answers to these questions depend on the spin structure of the hadron and the hadronization mechanism. They can even be different in the longitudinally polarized case from those for the transversely polarized case. Since neither question can be resolved using perturbative calculations at present, phenomenological studies are need to search for the answers. Currently, there exist two distinct pictures for the spin structure of the nucleon, i.e., the SU(6) picture based on the SU(6) wave function of the baryon, and the deeply inelastic scattering picture based on polarized deeply inelastic lepton-nucleon scattering data and other inputs such as symmetry assumptions and data from other experiments. It is of particular interest to know which one is suitable here.

It is clear that to study these questions, one needs to know the polarization of the quark before fragmentation and measure the polarization of the hadron produced in the fragmentation. Hence, we have the following two possibilities. One is to study hyperon polarization in e^+e^- annihilation at the Z^0 pole, in polarized *ep* deeply inelastic scattering or in high p_T polarized *pp* collisions. The other is to study the vector meson polarization in these processes.

It has been pointed out that the Λ polarization in e^+e^- annihilation at the Z^0 pole provides a very special test of the applicability of the SU(6) picture in the longitudinally polarized case. This is because, for $e^+e^-\to Z^0\to\Lambda X$, the $s\bar{s}$ created at the e^+e^- annihilation vertex is almost completely longitudinally polarized. In this case, if we assume that the quark keeps its polarization in the fragmentation and use the SU(6) wave function to connect the spin of the quark and that of the hyperon that contains this quark, we should obtain a maximum for the magnitude of Λ polarization since in this picture the spin of the *s* quark is completely transferred to Λ. Experimental data were obtained for $e^+e^-\to Z^0\to\Lambda X$ by the ALEPH and OPAL Collaborations at the CERN e^+e^- collider LEP [21,22]. We showed that if we assume that the q^0 keeps its polarization in fragmentation and the SU(6) wave function can be used in relating the polarization of q^0 and that of the produced hyperon which contains q^0, we obtain results that are in agreement with the data [19]. This result is rather encouraging. But the accuracy and abundance of the data are not enough to make a conclusive judgment. In particular, there is no direct measurement available at all in the transversely polarized case. We therefore made a systematic calculation for hyperon polarizations in all the different reactions [23]. There are also data for Λ polarization in deeply inelastic scattering [24]; they are also consistent with the results obtained using the SU(6) picture.

There are in addition data that provide information on vector meson polarization in high energy reactions. The 00 elements of the helicity density matrices for K^*, ρ, etc., in $e^+e^-\to Z^0\to VX$ have been measured by the ALEPH, DELPHI and OPAL Collaborations at LEP [31–33]. We showed that these data can also be understood using the SU(6) picture [34]. Further tests, not only in the longitudinally polarized case but also in the transversely polarized case, are under way. In this paper, we assume the picture is the same in the longitudinally and transversely polarized cases and use it in studying hyperon polarization in unpolarized high energy hadron-hadron collisions.

With the two points discussed in the last and this subsections, we can already obtain the polarizations for those hyperons that have one valence quark of the same flavor as that of the projectile, e.g., $pp\to\Sigma^-X$, $K^-p\to\Lambda X$, and $\Sigma^-p\to\Sigma^+X$. Some of the qualitative features of the results are given in Ref. [10]. It is encouraging to see that all of them are in good agreement with the data [3].

C. Correlation between the spin of q_s and that of $\bar{q}_s$ which combine with the (q_vq_v) and remaining q_v to form the produced hyperon and the associated meson

To study the polarization of hyperons such as Λ in $pp\to\Lambda X$, i.e., those that have two valence quarks with the same flavors as those of the projectile, we encounter the following question: If a hadron is produced by two valence quarks (a valence diquark) q_vq_v of the projectile, the remaining valence quark q_v produces an associated hadron. What are the spin states of the q_s and $\bar{q}_s$ that combine with the q_vq_v and the remaining q_v to form, respectively, the produced hyperon and the associatively produced meson in the fragmentation region?

Theoretically, it is quite difficult to derive this since we are in the very small *x* region; the production of such pairs is also of soft nature in general and cannot be calculated using perturbative theory. To get some clue to this problem, we still start from the single-spin left-right asymmetry A_N. The existing data [14,15] clearly show that A_N in $p^\uparrow p\to\Lambda X$ is large in magnitude and negative in sign in the fragmentation region. We note that the Λ in this region is mainly produced by the valence (u_vd_v) diquark of the projectile and is associated with the production of a K^+ by the remaining u_v. From the SU(6) wave function we learn that (u_vd_v) has to be in the spin zero state and the spin of the proton is carried by the remaining u_v. According to the $\vec{s}_q\cdot\vec{n}$ spin correlation discussed in Sec. II A, A_N for the associatively produced K^+ is positive. Hence, to understand the data on A_N for Λ, we simply need to assume that the Λ produced by (u_vd_v) and the K^+ that is associatively produced by the remaining u_v move in opposite transverse directions [35], which is just a direct consequence of transverse momentum conservation. Now, we apply this result to unpolarized *pp* collisions and consider the case that a Λ is produced by (u_vd_v) together with an *s* quark and a K^+ is associatively produced by the remaining u_v together with the $\bar{s}$. According to the $\vec{s}_q\cdot\vec{n}$ spin correlation mentioned in Sec. II A, the remaining u_v should have a large probability of being polarized in the $-\vec{n}$ direction since the normal to the production plane for the K^+ is opposite to the normal $\vec{n}$ to the production plane of the Λ. The polarization is $-C$. Since K^+ is a spin zero object, the $\bar{s}$ should have a polarization of $+C$ (in the $\vec{n}$ direction). The data for Λ polarization show that P_Λ is large and negative. This means that the *s* quark has a negative polarization. We thus reach the conclusion that, to understand the polarization of Λ in $pp\to\Lambda X$, we need to assume that the s and $\bar{s}$ have opposite spins. Under this assumption, together with the two points mentioned in the last two subsections, we obtained a good fit to the data on Λ polarization [10].

This point, of course, needs to be further studied and tested experimentally. We found that the simple exclusive process are most suitable for this purpose. Hyperon polarizations in exclusive processes such as $pp\to p\Lambda K^+$ and $e^-p\to e^-\Lambda K^+$ are very sensitive to the spin states of the s and $\bar{s}$ pairs. If the s and $\bar{s}$ have opposite spins, the results obtained for Λ polarization in $pp\to p\Lambda K^+$ should take the maximum among the different channels for $pp\to\Lambda X$. This is because here we have a situation where the Λ is definitely produced by the $(u_v d_v)$ valence diquark and is definitely associated with a K^+ that is definitely produced by the remaining u_v of the incident p. Hence, we obtain in this case $P_\Lambda=-C=-0.6$. This is in good agreement with the data obtained by the R608 Collaboration at CERN which show that $P_\Lambda=-0.62\pm0.04$ [17].

We also calculated the Λ polarization in $e^-p\to e^-\Lambda K^+$ in all three possible spin states of s and $\bar{s}$, i.e., opposite, the same, or uncorrelated [16]. We found that the results in the three cases are quite different from each other. Now, experimental data have been obtained for Λ polarization in $e^-p\to e^-\Lambda K^+$ by the CLAS Collaboration at Jefferson Laboratory [18]. Comparing to the data, we see that the results obtained in the case that the spins of s and $\bar{s}$ are opposite are favored. Further tests are also under way.

We now assume that this is true in general, i.e., the q_s that combines with the $(q_v q_v)$ of the projectile to form the hyperon and the $\bar{q}_s$ that combines with the remaining q_v to form the associatively produced meson have opposite spins. Under this assumption, we obtain the result that the q_s should be polarized in the $-\vec{n}$ direction and the polarization is $-C$. We apply this to the production of different hyperons to calculate the hyperon polarizations in unpolarized high energy hadron-hadron or hadron-nucleus collisions in the next section.

We emphasize that the above mentioned result is true for the production of hyperons associated with the production of pseudoscalar mesons. The situation should be different if the associatively produced meson is a vector meson. This influence will be further investigated in a separate paper [36] and here we consider only the former case.

III. THE CALCULATION METHOD

Using the picture discussed in the last section, we can calculate the hyperon polarization in different hadron-hadron collisions. We now present the formulas used in these calculations.

A. General formulas

We consider the process $P+T\to H_i+X$, where P and T denote, respectively, the projectile and target hadrons, and H_i denotes the ith kind of hyperon. The hyperon polarization P_{H_i} is defined as

$$P_{H_i}(x_F|s)\equiv\frac{N(x_F,H_i,\uparrow|s)-N(x_F,H_i,\downarrow|s)}{N(x_F,H_i,\uparrow|s)+N(x_F,H_i,\downarrow|s)}=\frac{\Delta N(x_F,H_i|s)}{N(x_F,H_i|s)},\tag{3}$$

where $N(x_F,H_i,l|s)$ is the number density of H_i's polarized in the same $(l=\uparrow)$ or opposite $(l=\downarrow)$ direction as the normal $(\vec{n})$ to the production plane at a given $\sqrt{s}$; $x_F\equiv2p_\parallel/\sqrt{s}$, $p_\parallel$ is the longitudinal momentum of H_i with respect to the incident direction of P, and $\sqrt{s}$ is the total c.m. energy of the colliding hadron system. It is clear that the denominator is nothing else but the number density of H_i without specifying the polarization.

To calculate $\Delta N(x_F,H_i|s)$, we divide the final hyperons H_i into the following four groups according to the different origins for the production: (A) those directly produced and containing a valence diquark (two valence quarks) $(q_v q_v)^P$ of the projectile; (B) those directly produced and containing a valence quark q_v^P of the projectile; (C) those from the decay of the directly produced heavier hyperons H_j that contain a $(q_v q_v)^P$ or a q_v^P; and (D) the others. In this way, we have

$$N(x_F,H_i|s)=N_0(x_F,H_i|s)+D^A(x_F,H_i|s)+\sum_f D^{B,f}(x_F,H_i|s)+\sum_j D^{C,H_j}(x_F,H_i|s),\tag{4}$$

where $D^A(x_F,H_i|s)$, $D^{B,f}(x_F,H_i|s)$, and $D^{C,H_j}(x_F,H_i|s)$ denote the contributions from groups A, B, and C, respectively; the superscripts f and j denote the flavor of q_v and the type of H_j, respectively; N_0 is the contribution from D.

According to the picture discussed in the last section, hyperons from groups A, B, and C can be polarized, while those from D are not. This means that those from A, B, and C contribute to the numerator ΔN of Eq. (4), i.e., we have

$$\Delta N(x_F,H_i|s)=\Delta D^A(x_F,H_i|s)+\sum_f \Delta D^{B,f}(x_F,H_i|s)+\sum_j \Delta D^{C,H_j}(x_F,H_i|s).\tag{5}$$

Since valence quarks usually carry large fractions of the momenta of the incident hadrons, we expect that, for very large x_F, $D^A(x_F,H_i|s)$ dominates. For small x_F, N_0 dominates, while for moderate x_F, $D^B(x_F,H_i|s)$ plays the dominant role. Hence, if we neglect the decay contributions, we expect from Eqs. (3)–(5) that $P_{H_i}(x_F|s)$ has the following general properties. For x_F increasing from 0 to 1, it starts from 0, increases to $\Sigma_f\Delta D^{B,f}(x_F,H_i|s)/\Sigma_f D^{B,f}(x_F,H_i|s)$, and finally tends to $\Delta D^A(x_F,H_i|s)/D^A(x_F,H_i|s)$ at $x_F\to1$. We will come to this point in the next section for particular hyperons in the specified reaction. We now first discuss the calculations of all these D's, ΔD's, and N_0 in the following.

B. Calculations of *D*'s and N_0

The contributions of hyperons from the different groups discussed above are entirely determined by the hadronization mechanisms in the unpolarized case. They are independent of the polarization of the hadrons. We can calculate them using a hadronization model that gives a good description of the unpolarized data. For this purpose, the simple model used in Refs. [25,26,35] is a very practical choice. In this model, hyperons from groups A and B are described as the products of the following "direct-formation" or "direct-fusion" process. For A, it is

$$(q_v q_v)^P + q_s^T \to H_i ,$$

and for B, it is

$$q_v^P + (q_s q_s)^T \to H_i ,$$

where q_s^T and $(q_s q_s)^T$ denote a sea quark or a sea diquark from the target. The number densities of the hyperons produced in these processes are determined by the number densities of the initial partons. They are given by [25,26,35]

$$D^A(x_F, H_i|s) = \kappa^d_{H_i} f_D^P(x^P|q_v q_v) q_s^T(x^T), \tag{6}$$

$$D^{B,f}(x_F, H_i|s) = \kappa_{H_i} q_v^P(x^P) f_D^T(x^T|q_s q_s), \tag{7}$$

where $x^P \approx x_F$ and $x^T \approx m^2_{H_i}/(s x_F)$, following from energy-momentum conservation in the direct formation processes; $q_i(x)$ is the quark distribution function, where q denotes the flavor of the quark and the subscript $i = v$ or s denotes whether it is for valence or sea quarks; $f_D(x|q_i q_j)$ is the diquark distribution function, where $q_i q_j$ denotes the flavor and whether they are valence or sea quarks, and the superscripts P or T denote the name of the hadron; $\kappa^d_{H_i}$ and κ_{H_i} are two constants which are fixed by fitting two data points in the large x_F region.

Since most of the decay processes $H_j \to H_i + X$ that we consider are two-body decays, $D^{C,H_j}(x_F, H_i|s)$ can be calculated from a convolution of $D^A(x_F, H_j|s)$ or $D^{B,f}(x_F, H_j|s)$ with the distribution describing the decay process. The calculations are in principle straightforward, but in practice a little complicated, and detailed information about the transverse momentum distribution of H_j is needed. Since the influence is not very large, we use the following approximation for simplicity. We neglect the distribution caused by the decay process and take the average value for x_F instead. More precisely, we take $H_j \to H_i + M$ as an example (where M denotes a meson). For an H_j with a given longitudinal momentum fraction x_F^j, the resulting x_F of the produced H_i can take different values. The distribution of x_F at a fixed x_F^j can be obtained from the isotropic distribution of the momenta of the decay products in the rest frame of H_j. This can be calculated if the transverse momentum of H_j is also given. The average value of the resulting x_F has a simple expression, $\langle x_F(H_j \to H_i)\rangle = x_F^j E_{H_i,H_j}/m_{H_j}$, where $E_{H_i,H_j} = (m^2_{H_j} + m^2_{H_i} - m^2_M)/(2 m_{H_j})$ is the energy of H_i in the rest frame of H_j. We see that $\langle x_F(H_j \to H_i)\rangle$ is independent of the transverse momentum of H_j. In our calculations, we simply neglect the distribution and take $x_F = \langle x_F(H_j \to H_i)\rangle$ for a given x_F^j. In this approximation, we have

$$D^{C,H_j}(x_F, H_i|s) \approx Br(H_j \to H_i) \frac{m_{H_j}}{E_{H_i,H_j}} \left[D^A\left(\frac{m_{H_j} x_F}{E_{H_i,H_j}}, H_j|s \right) + \sum_f D^{B,f}\left(\frac{m_{H_j} x_F}{E_{H_i,H_j}}, H_j|s \right) \right], \tag{8}$$

where $Br(H_j \to H_i)$ is the branch ratio for the decay channel.

Having calculated all these D's, we can obtain the N_0 by parametrizing the difference of the experimental data using the number density $N(x_F, H_i|s)$ of produced H_i and these D's. We emphasize that the direct-fusion model has been proposed to describe the production of hadrons in the fragmentation region. As has been shown by comparing different parts of the contributions to N in Refs. [25,26,35], this mechanism plays the dominating role in the large x_F region such as $x_F \gtrsim 0.5$–0.6. It is clear that nobody knows *a priori* whether the quark distribution functions can be used in describing the quark-fusion process which leads to the hadrons in the fragmentation region with moderately large p_T. The applicability follows from the empirical facts pointed out by Ochs [37,38], and the phenomenological work by Das and Hwa [39] a long time ago. It has been pointed out that various experiments have shown that the longitudinal momentum distributions of the produced hadrons in the fragmentation region are very similar to those of the corresponding valence quarks in the colliding hadrons [37,38]. The model follows directly from this observation. It is interesting to note that this simple model not only is consistent with the observation of Ochs [37,38] already made in 1977 and the theoretical work by Das and Hwa [39], but also has experienced a number of tests such as isospin invariance of N_0, etc. (for a summary, see Ref. [25]). Furthermore, the energy dependence that is contained in D due to the energy dependence of x^T leads [40] naturally to the energy dependence of the single-spin left-right asymmetry A_N observed by the BNL E925 Collaboration [41] compared with those by the Fermilab E704 Collaboration. We use this model to calculate the D's and N_0 in Eq. (4). The quark distributions are taken from parametrization at low Q^2 such as $Q \approx 1$ GeV/c.

C. Calculations of ΔD's

The calculations of the differences, i.e., ΔD^A, $\Delta D^{B,f}$, and $\Delta D^{C,H_j}$, are the core ingredients of this model. They are described in the following.

1. Calculation of ΔD^A

This is to determine the polarization of H_i coming from group A, i.e., $(q_v q_v)^P + q_s^T \to H_i$. Here, as mentioned in the last section, we consider only the case that $(q_v q_v)^P + q_s^T \to H_i$ is associated with the production of $q_v + \overline{q_s} \to M$ where M is a pseudoscalar meson. We recall that, according to the third point discussed in Sec. II, q_s should be polarized in the

$-\vec{n}$ direction and the polarization is $-C$. Hence, to determine the polarization of H_i, we need to know the relative weights for $(q_v q_v)^P$ to be in the different spin states $|(q_v q_v)_{s_d,s_{dn}}\rangle$ where the subscripts s_d and s_{dn} denote the spin and its n component of the diquark $q_v q_v$. Since the $q_v q_v$ is from the projectile P, these relative weights can be calculated using the SU(6) wave function of P. In this way, we obtain the relative weights for the production of $(q_v q_v) q_s$ in different spin states $|(q_v q_v)_{s_d,s_{dn}} q_s^{\downarrow}\rangle$. We denote this relative weight by $w(s_d,s_{dn}|q_v q_v)$. After that, we make the projections of these different spin states $|(q_v q_v)_{s_d,s_{dn}} q_s^{\downarrow}\rangle$ to the wave functions $|H_i(s_n)\rangle$ of H_i with different values of s_n (which denotes the projection of the spin of H_i along the $\vec{n}$ direction) and obtain the relative weights for H_i to be in different spin states. The polarization of such H_i is then given by

$$P_{H_i}^A=\frac{\sum_{s_n,s_d,s_{dn}} w(s_d,s_{dn}|q_v q_v)\cdot|\langle (q_v q_v)_{s_d,s_{dn}} q_s^{\downarrow}|H_i(s_n)\rangle|^2\cdot s_n}{\sum_{s_n,s_d,s_{dn}} w(s_d,s_{dn}|q_v q_v)\cdot|\langle (q_v q_v)_{s_d,s_{dn}} q_s^{\downarrow}|H_i(s_n)\rangle|^2\cdot s_{n,max}} C\equiv \alpha_{H_i}^A C. \tag{9}$$

Hence, the difference ΔD^A is given by

$$\Delta D^A(x_F,H_i|s)=\alpha_{H_i}^A C D^A(x_F,H_i|s). \tag{10}$$

For different reactions, the corresponding $\alpha_{H_i}^A$'s are calculated and the results are presented in the next section.

2. Calculation of $\Delta D^{B,f}$

The polarization of H_i coming from group B, i.e., $q_v^P + (q_s q_s)^T \to H_i$, is determined in the following way.

First, using the first point discussed in Sec. II, we determine the polarization of q_v. As given by Eq. (2), it is polarized in the $\vec{n}$ direction and the polarization is C. Second, the corresponding relative probabilities for all possible spin states of $(q_s q_s)$ are taken as the same. This means that $q_v(q_s q_s)$ has equal probability of 1/4 to be in the different spin states $|q_v^{\uparrow}(q_s q_s)_{s_d,s_{dn}}\rangle$ where $s_d=0,1$ and s_{dn} takes all possible values. Finally, we project the different spin states of $q_v(q_s q_s)$ to $|H_i(s_n)\rangle$ to calculate the relative weights for H_i to be in the different spin states. Then the polarization of such H_i is given by

$$P_{H_i}^{B,f}=\frac{\sum_{s_n,s_d,s_{dn}} |\langle q_v^{\uparrow}(q_s q_s)_{s_d,s_{dn}}|H_i(s_n)\rangle|^2\cdot s_n}{\sum_{s_n,s_d,s_{dn}} |\langle q_v^{\uparrow}(q_s q_s)_{s_d,s_{dn}}|H_i(s_n)\rangle|^2\cdot s_{n,max}} C\equiv \alpha_{H_i}^{B,f} C. \tag{11}$$

We note that, for a spin 1/2 hyperon H_i, we have

$$\alpha_{H_i}^{B,f}=2\langle s_{f,n}\rangle, \tag{12}$$

which is nothing else but the polarization of the valence quark q of flavor f in $H_i^{\uparrow}$ along the polarization of H_i. It is just the fragmentation spin transfer factor $t_{H_i,f}^F$, which is defined as the probability for the polarization of q_v to be transferred to H_i in the fragmentation $q_v \to H_i + X$ in the case that the q_v is contained in H_i [23]. Obviously, using the SU(6) wave function of H_i, we can calculate these $t_{H_i,f}^F$'s for different H_i and quark flavor f. The results are given in Table I. Hence, the difference $\Delta D^{B,f}$ is given by

$$\Delta D^{B,f}(x_F,H_i|s)=C t_{H_i,f}^F D^{B,f}(x_F,H_i|s). \tag{13}$$

3. Calculation of $\Delta D^{C,H_j}$

To determine the polarization of H_i from the decay process $H_j \to H_i + X$, we should first determine the polarization of H_j. Since H_j is directly produced and the origins belong to group A or B discussed above, its polarization is determined in the same way as in the last two cases. The polarization of H_j can be transferred to H_i in the decay process $H_j \to H_i + X$. The probability is denoted by $t_{i,j}^D$ and is called the decay spin transfer factor. This means that

$$\Delta D^{C,H_j}(x_F,H_i|s)\approx C t_{i,j}^D Br(H_j\to H_i) \times \frac{m_{H_j}}{E_{H_i,H_j}}\left[\alpha_{H_j}^A D^A\left(\frac{m_{H_j}x_F}{E_{H_i,H_j}},H_j|s\right) + \sum_f t_{H_j,f}^F D^{B,f}\left(\frac{m_{H_j}x_F}{E_{H_i,H_j}},H_j|s\right)\right]. \tag{14}$$

The decay spin transfer factor $t_{i,j}^D$ is universal in the sense that it is determined by the decay process and is independent of the process by which H_j is produced. For $\Sigma^0 \to \Lambda\gamma$, it is

TABLE I. Fragmentation spin transfer factor $t_{H_i,f}^F$ obtained using the SU(6) wave function.

	Λ	Σ^+	Σ^0	Σ^-	Ξ^0	Ξ^-
u	0	2/3	2/3	0	$-1/3$	0
d	0	0	2/3	2/3	0	$-1/3$
s	1	$-1/3$	$-1/3$	$-1/3$	2/3	2/3

determined as $t^D_{\Lambda,\Sigma^0}=-1/3$ [42]; and for $\Xi\to\Lambda\pi$, $t^D_{\Lambda,\Xi}=(1+2\gamma)/3$, where $\gamma=0.87$ is a constant and can be found in the "Review of Particle Properties" [43].

IV. RESULTS AND DISCUSSIONS

In this section, we apply the calculation method given in last section to hyperons in different reactions, and present the results obtained.

A. Hyperon polarization in $p+p\to H_i+X$

We first consider *pp* or *pA* collisions and present the results for different hyperons. As we can see from the last section, the hyperon polarizations in the proposed picture depend on the valence quarks of the projectile and the sea of the target. Hence, if we neglect the small influence from the differences between the sea of the proton and that of the neutron, we should obtain the same results for *pp*, *pn*, or *pA* collisions.

1. $p+p\to\Lambda+X$

The calculations of P_Λ in $pp\to\Lambda X$ have been given in Ref. [10]. For completeness, we summarize the results here.

For Λ production in *pp* collisions, there is one contributing process from group A, i.e., $(u_vd_v)^P+s_s^T\to\Lambda$. There are two contributing processes from B, i.e., $u_v^P+(d_ss_s)^T\to\Lambda$ and $d_v^P+(u_ss_s)^T\to\Lambda$. Thus, we have

$$D^A(x_F,\Lambda|s)=\kappa^d_\Lambda f_D(x^P|u_vd_v)s_s(x^T), \tag{15}$$

$$D^{B,u}(x_F,\Lambda|s)=\kappa_\Lambda u_v(x^P)f_D(x^T|d_ss_s), \tag{16}$$

$$D^{B,d}(x_F,\Lambda|s)=\kappa_\Lambda d_v(x^P)f_D(x^T|u_ss_s), \tag{17}$$

where $x^P\approx x_F$ and $x^T\approx m^2_\Lambda/(sx_F)$, following from energy-momentum conservation in the direct-formation processes; κ^d_Λ and κ_Λ are two constants. Here, as usual, we omit the superscripts of the distribution functions when they are for protons.

We take the contributions of $J^P=(1/2)^+$ hyperon decay into account. For Λ, we have $\Sigma^0\to\Lambda\gamma$ and $\Xi^{0,-}\to\Lambda\pi^{0,-}$. For $pp\to\Sigma^0X$, we have similar contributing processes from groups A and B as those for $pp\to\Lambda X$, i.e., $(u_vd_v)^P+s_s^T\to\Sigma^0$, $u_v^P+(d_ss_s)^T\to\Sigma^0$, and $d_v^P+(u_ss_s)^T\to\Sigma^0$. The number densities of Σ^0 from these processes are given by

$$D^A(x_F,\Sigma^0|s)=\kappa^d_{\Sigma^0}f_D(x^P|u_vd_v)s_s(x^T), \tag{18}$$

$$D^{B,u}(x_F,\Sigma^0|s)=\kappa_\Sigma u_v(x^P)f_D(x^T|d_ss_s), \tag{19}$$

$$D^{B,d}(x_F,\Sigma^0|s)=\kappa_\Sigma d_v(x^P)f_D(x^T|u_ss_s), \tag{20}$$

respectively. Here, $x^P\approx x_F$ and $x^T\approx m^2_\Sigma/(sx_F)$, $\kappa^d_{\Sigma^0}$ and κ_Σ are two corresponding constants for Σ^0 production. For $pp\to\Xi^0X$ and $pp\to\Xi^-X$, we have no contributing process from A but some from B. They are $u_v^P+(s_ss_s)^T\to\Xi^0$ and $d_v^P+(s_ss_s)^T\to\Xi^-$, respectively. The corresponding number densities are given by

$$D^{B,u}(x_F,\Xi^0|s)=\kappa_\Xi u_v(x^P)f_D(x^T|s_ss_s), \tag{21}$$

$$D^{B,d}(x_F,\Xi^-|s)=\kappa_\Xi d_v(x^P)f_D(x^T|s_ss_s), \tag{22}$$

where $x^P\approx x_F$ and $x^T\approx m^2_\Xi/(sx_F)$.

As discussed in Sec. III, their contributions to Λ production are given by

$$D^{C,\Sigma^0}(x_F,\Lambda|s)\approx\frac{m_\Sigma}{E_{\Lambda,\Sigma^0}}\left[D^A\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)+\sum_{f=u,d}D^{B,f}\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)\right], \tag{23}$$

$$D^{C,\Xi^0}(x_F,\Lambda|s)\approx\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,u}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^0|s\right), \tag{24}$$

$$D^{C,\Xi^-}(x_F,\Lambda|s)\approx\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,d}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right), \tag{25}$$

respectively. Hence, we obtain

$$N(x_F,\Lambda|s)=N_0(x_F,\Lambda|s)+D^A(x_F,\Lambda|s)+\sum_{f=u,d}D^{B,f}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}D^{C,H_j}(x_F,\Lambda|s). \tag{26}$$

The different weights $w(s_d,s_{dn}|u_vd_v)$ for (u_vd_v) from protons in different spin states can be calculated by rewriting the SU(6) wave function of the proton as follows:

$$|p^\uparrow\rangle=\frac{1}{2\sqrt{3}}[3u^\uparrow(ud)_{0,0}+u^\uparrow(ud)_{1,0}-\sqrt{2}u^\downarrow(ud)_{1,1}]. \tag{27}$$

Taking into account that the proton is unpolarized and thus has equal probabilities to be in $|p^\uparrow\rangle$ and $|p^\downarrow\rangle$, we obtain the results for $w(s_d,s_{dn}|u_vd_v)$ as shown in Table II. We then project the different spin states of $(u_vd_v)s_s$ on the wave function of Λ and obtain the relative weights for the production of Λ from this process in different spin states in Table II. From these results, we obtain also the α^A_Λ as shown in the table. Hence its contribution to $\Delta N(x_F,\Lambda|s)$, i.e., $\Delta D^A(x_F,\Lambda|s)=C\alpha^A_\Lambda D^A(x_F,\Lambda|s)$, can also be calculated. Similarly, we obtain the corresponding results for Σ^0 production and show them in the same table. From these results, we obtain that

$$\Delta N(x_F,\Lambda|s)=C\alpha^A_\Lambda D^A(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}\Delta D^{C,H_j}(x_F,\Lambda|s), \tag{28}$$

TABLE II. Relative weights $w(s_d,s_{dn}|u_v d_v)$ for the $u_v d_v$ diquark from the proton in different spin states, those for the produced Λ and Σ^0, the corresponding α_Λ^A and $\alpha_{\Sigma^0}^A$, and the total weights w_Λ^A and $w_{\Sigma^0}^A$.

Possible spin states	$(u_v d_v)_{0,0}s_s^\downarrow$		$(u_v d_v)_{1,0}s_s^\downarrow$		$(u_v d_v)_{1,1}s_s^\downarrow$	
$w(s_d,s_{dn}\|u_v d_v)$	3/4		1/12		1/6	
Possible products	$\Lambda^\downarrow$	$\Sigma^{0\downarrow}$	$\Lambda^\downarrow$	$\Sigma^{0\downarrow}$	$\Lambda^\uparrow$	$\Sigma^{0\uparrow}$
$\|\langle(q_v q_v)_{s_d,s_{dn}}q_s^\downarrow\|H_i(s_n)\rangle\|^2$	1	0	0	1/3	0	2/3
The final relative weights	3/4	0	0	1/36	0	1/9
The resulting w_{Hi}^A and α_{Hi}^A	Λ: 3/4, -1;				Σ^0: 5/36, 3/5	

$$\Delta D^{C,\Sigma^0}(x_F,\Lambda|s)\approx Ct^D_{\Lambda,\Sigma^0}\frac{m_\Sigma}{E_{\Lambda,\Sigma^0}}\left[\alpha^A_{\Sigma^0}D^A\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)+\sum_{f=u,d}t^F_{\Sigma^0,f}D^{B,f}\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)\right], \tag{29}$$

$$\Delta D^{C,\Xi^0}(x_F,\Lambda|s)\approx Ct^D_{\Lambda,\Xi}t^F_{\Xi^0,u}\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,u}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^0|s\right), \tag{30}$$

$$\Delta D^{C,\Xi^-}(x_F,\Lambda|s)\approx Ct^D_{\Lambda,\Xi}t^F_{\Xi^-,d}\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,d}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right), \tag{31}$$

where the fragmentation spin transfer factors $t^F_{H_j,f}$ for different quark flavors f and different hyperons H_j are listed in Table I, and the decay spin transfer factors t^D_{Λ,H_j} are given in Sec. III.

Using the results given by Eqs. (15)–(31), we can now calculate $P_\Lambda(x_F|s)$ as a function of x_F. Before we show the numerical results, we first look at the qualitative features. Just as mentioned at the end of Sec. III A for $pp\to\Lambda X$, we expect that $D^A(x_F,\Lambda|s)$, i.e., the contribution from $(u_v d_v)^P+s_s^T\to\Lambda$, dominates $N(x_F,\Lambda|s)$ at $x_F\to1$. $N_0(x_F,\Lambda|s)$ dominates for very small x_F, while $D^{B,u}(x_F,\Lambda|s)$ and $D^{B,d}(x_F,\Lambda|s)$, i.e., those from $u_v^P+(d_s s_s)^T\to\Lambda$ and $d_v^P+(u_s s_s)^T\to\Lambda$, play the dominating role for moderate x_F. Since $\alpha^A_\Lambda=-1$ and $t^F_{\Lambda,u}=t^F_{\Lambda,d}=0$, we expect that, if the decay contribution can be neglected, for x_F going from 0 to 1, $P_\Lambda(x_F|s)$ starts from 0, becomes nonzero quite slowly, and tends to $C\alpha^A_\Lambda=-C$ at $x_F=1$. Taking the decay contribution into account, we expect a significant contribution from $\Sigma^0\to\Lambda\gamma$ which leads to a negative $P_\Lambda(x_F|s)$ at moderate x_F and makes the $|P_\Lambda(x_F|s)|$ less than C at x_F near 1.

We now use Eqs. (15)–(31) to get the numerical results for $P_\Lambda(x_F|s)$ as a function of x_F. The only unknown parameters are the κ's, which should be determined by the unpolarized experimental data on the number densities for the corresponding hyperons. To reduce the arbitrariness in determining these κ's, we take the same κ_{H_i}'s for different hyperons produced in processes of group B. For those from group A, we take the $\kappa^d_{H_i}$'s for different hyperons as

$$\kappa^d_{H_i}=w^d_{H_i}\cdot\kappa^d, \tag{32}$$

where κ^d is taken as a constant independent of H_i, and

$$w^d_{H_i}\equiv\sum_{s_d,s_{dn},s_n}w(s_d,s_{dn}|q_v q_v)\cdot|\langle(q_v q_v)_{s_d,s_{dn}}q_s^\downarrow|H_i(s_n)\rangle|^2 \tag{33}$$

is the total relative weight for the production of H_i from $(q_v q_v)^P+q_s^T\to H_i$. This means that only the H_i dependence of $\kappa^d_{H_i}$ from the spin statistics is taken into account. The $w^d_{H_i}$'s for Λ and Σ^0 are given in Table II. In this way, we have only two free κ's, i.e., κ^d for $(q_v q_v)^P+q_s^T\to H_i$ and $\kappa_{Hi}\equiv\kappa$ for $q_v^P+(q_s q_s)^T\to H_i$. We determine them by using the data for $N(x_F,\Lambda|s)$ [44].

Having the κ's, we calculate $P_\Lambda(x_F|s)$ as a function of x_F. For the unpolarized quark distribution functions, we use the Glück-Reya-Vogt (GRV) 98 leading order (LO)set [45]. For the unpolarized diquark distribution $f_D(x^P|q_v q_v)$, we use the parametrization given in Ref. [46]. For the unpolarized diquark distribution $f_D(x^T|q_s q_s)$, we simply use a convolution of $q_s(x)$ for the two sea quarks. The results obtained for $P_\Lambda(x_F|s)$ are given in Fig. 1. The data in the figure are only those for $p_T>1$ GeV/c, because the C that we used in the calculations was determined by A_N for the p_T interval $0.7<p_T<2.0$ GeV/c. We see clearly that, as x_F changes from 0 to 1, the $P_\Lambda(x_F|s)$ obtained indeed starts from 0, goes slowly to about -15%, and finally to about -50% at x_F

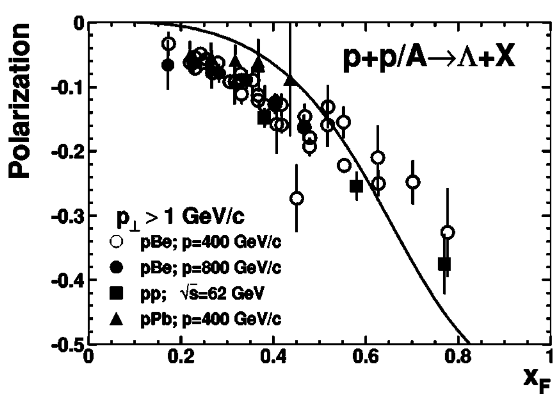

FIG. 1. Calculated results for the polarization of Λ in pp collisions as a function of x_F. Data are taken from Refs. [47–49].

TABLE III. Relative weights $w(s_d,s_{dn}|u_v u_v)$ for the $u_v u_v$ diquark from protons in different spin states, those for the produced Σ^+, the corresponding $\alpha^A_{\Sigma^+}$, and the total relative weight $w^A_{\Sigma^+}$.

Possible spin states	$(u_v u_v)_{1,0} s_s^{\downarrow}$	$(u_v u_v)_{1,1} s_s^{\downarrow}$
$w(s_d,s_{dn}\|u_v u_v)$	1/3	2/3
Possible products	$\Sigma^{+\downarrow}$	$\Sigma^{+\uparrow}$
$\|\langle (u_v u_v)_{s_d,s_{dn}} s_s^{\downarrow}\|\Sigma^+(s_n)\rangle\|^2$	1/3	2/3
The final relative weights	1/9	4/9
The resulting $w^A_{\Sigma^+}$ and $\alpha^A_{\Sigma^+}$	$w^A_{\Sigma^+}=5/9,\ \alpha^A_{\Sigma^+}=3/5$	

$\rightarrow 1$. These qualitative features are the same as we expect from the qualitative analysis above and are in good agreement with the data [47–49].

It should be mentioned that, in the calculations presented above, we considered only the associated production of hyperons of group A, i.e., $(q_v q_v)^P + q_s^T \rightarrow H_i$, with a pseudoscalar meson M, i.e., $q_v^P + \overline{q_s^T} \rightarrow M$. It is clear that the associatively produced hadron can also be a vector meson. Taking this into account, we expect that the correlations between the spins of these quarks will be reduced. As a consequence, the polarizations of hyperons from group A should be slightly reduced. Since such mesons contribute mainly at large x_F, this effect should make $|P_\Lambda|$ smaller than those presented in Fig. 1 in the large x_F region. From the figure, we also see that there is indeed room left for such an effect. Similar effects exist for $pp \rightarrow \Sigma^+ X$ and $\Sigma^- p \rightarrow \Sigma^-(\text{or } \Xi^-)X$ which will be discussed in the following. How large this effect can be is determined by the hadronization mechanisms. Presently, we are working on an estimation. The results will be published separately [36].

2. $p+p \rightarrow \Sigma^{\pm}$(or $\Xi^{0,-}$)$+X$

In a completely similar way, we calculate hyperon polarizations for $pp \rightarrow \Sigma^{\pm}$ or $\Xi^{0,-})X$. The results are given in the following.

For $pp \rightarrow \Sigma^+ X$, there is one contributing process from group A, i.e., $(u_v u_v)^P + s_s^T \rightarrow \Sigma^+$, and one from group B, i.e., $u_v^P + (u_s s_s)^T \rightarrow \Sigma^+$. The corresponding number densities are given by

$$D^A(x_F,\Sigma^+|s) = \kappa^d_{\Sigma^+} f_D(x^P|u_v u_v) s_s(x^T), \tag{34}$$

$$D^{B,u}(x_F,\Sigma^+|s) = \kappa u_v(x^P) f_D(x^T|u_s s_s), \tag{35}$$

where $x^P \approx x_F$ and $x^T \approx m_\Sigma^2/(s x_F)$. In the case that only $J^P = (1/2)^+$ hyperon decay is taken into account, there is no contribution from hyperon decay to Σ^+ production. Hence, we have

$$N(x_F,\Sigma^+|s) = N_0(x_F,\Sigma^+|s) + D^A(x_F,\Sigma^+|s) + D^{B,u}(x_F,\Sigma^+|s), \tag{36}$$

$$\Delta N(x_F,\Sigma^+|s) = C[\alpha^A_{\Sigma^+} D^A(x_F,\Sigma^+|s) + t^F_{\Sigma^+,u} D^{B,u}(x_F,\Sigma^+|s)], \tag{37}$$

where $t^F_{\Sigma^+,u}$ is given in Table I. $\alpha^A_{\Sigma^+}$ is calculated in completely the same way as α^A_Λ in the last subsection. The results are given in Table III. Thus, we have

$$P_{\Sigma^+}(x_F|s) = C\frac{\alpha^A_{\Sigma^+} D^A(x_F,\Sigma^+|s) + t^F_{\Sigma^+,u} D^{B,u}(x_F,\Sigma^+|s)}{N_0(x_F,\Sigma^+|s) + D^A(x_F,\Sigma^+|s) + D^{B,u}(x_F,\Sigma^+|s)}. \tag{38}$$

The situations for the production of Σ^-, Ξ^0, and Ξ^- in pp collisions are even simpler: There is no contributing process from group A and only one contributing process from B. When only $J^P = (1/2)^+$ hyperon decay is taken into account, there is also no contribution from hyperon decay to these hyperons. Hence, we have

$$P_{\Sigma^-}(x_F|s) = C\frac{t^F_{\Sigma^-,d} D^{B,d}(x_F,\Sigma^-|s)}{N_0(x_F,\Sigma^-|s) + D^{B,d}(x_F,\Sigma^-|s)}, \tag{39}$$

$$P_{\Xi^0}(x_F|s) = C\frac{t^F_{\Xi^0,u} D^{B,u}(x_F,\Xi^0|s)}{N_0(x_F,\Xi^0|s) + D^{B,u}(x_F,\Xi^0|s)}, \tag{40}$$

$$P_{\Xi^-}(x_F|s) = C\frac{t^F_{\Xi^-,d} D^{B,d}(x_F,\Xi^-|s)}{N_0(x_F,\Xi^-|s) + D^{B,d}(x_F,\Xi^-|s)}, \tag{41}$$

where the $t^F_{H_i,f}$'s are given in Table I;

$$D^{B,d}(x_F,\Sigma^-|s) = \kappa d_v(x^P) f_D(x^T|d_s s_s), \tag{42}$$

and $D^{B,u}(x_F,\Xi^0|s)$, and $D^{B,d}(x_F,\Xi^-|s)$ are given by Eqs. (21) and (22).

From these equations, we can calculate $P_{\Sigma^\pm}(x_F|s)$ and $P_\Xi(x_F|s)$ in pp collisions. Just as we mentioned at the end of Sec. III A that for $pp \rightarrow \Sigma^+ X$ we expect that $D^A(x_F,\Sigma^+|s)$ dominates at large x_F, $D^{B,u}(x_F,\Sigma^+|s)$ plays the dominating role at moderate x_F. Hence, we expect from Eq. (38) that $P_{\Sigma^+}(x_F|s) = 0$ for x_F near 0, increases to $Ct^F_{\Sigma^+,u} = 0.4$, and finally tends to $C\alpha^A_{\Sigma^+} = 0.36$ with increasing x_F. The situations for Σ^-, Ξ^0, and Ξ^- are even simpler since there is no contributing process from group A. Here, N_0 dominates at small x_F, while D^B dominates at large x_F. We thus expect that $P_{\Sigma^-}(x_F|s) = 0$ for x_F near 0 and increases to $Ct^F_{\Sigma^-,d}$

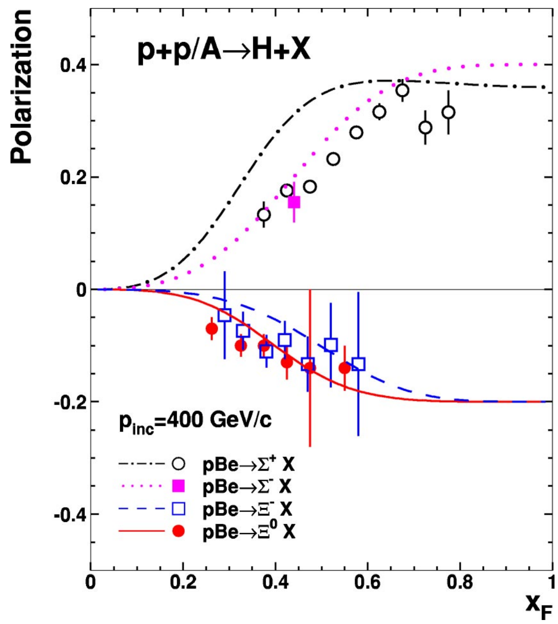

FIG. 2. Calculated results for the polarizations of $\Sigma^{\pm}$ and $\Xi^{0,-}$ at $p_{inc}=400$ GeV/c in pp collisions as functions of x_F. Data are taken from Refs. [50–53].

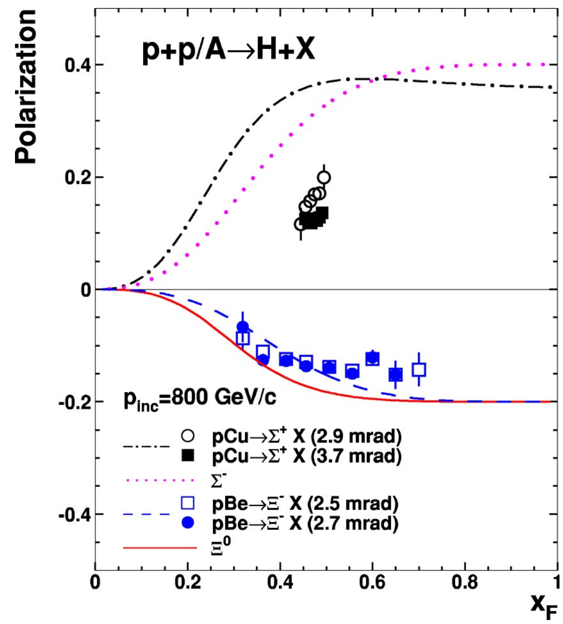

FIG. 3. Calculated results for the polarizations of $\Sigma^{\pm}$ and $\Xi^{0,-}$ at $p_{inc}=800$ GeV/c in pp collisions as functions of x_F. Data are taken from Refs. [54,55].

$=0.4$ with increasing x_F. For Ξ, $P_{\Xi}(x_F|s)=0$ for x_F near 0 and increases to $Ct^F_{\Xi,f}=-0.2$ with increasing x_F. To summarize, we expect that both $P_{\Sigma^+}(x_F|s)$ and $P_{\Sigma^-}(x_F|s)$ are positive in sign and increase fast to about 40% with increasing x_F. On the other hand, both $P_{\Xi^0}(x_F|s)$ and $P_{\Xi^-}(x_F|s)$ are negative in sign, and they decrease to -20% with increasing x_F. The magnitudes of $P_{\Sigma^{\pm}}$ are expected to be larger than those of P_{Ξ}. These qualitative features are in agreement with the available data [50–55].

We now use Eqs. (38)–(41) to obtain the numerical results for $P_{\Sigma^{\pm}}(x_F|s)$ and $P_{\Xi}(x_F|s)$ in pp collisions. The only unknown is N_0 for the corresponding hyperon. As we mentioned earlier in Sec. III, N_0 is independent of the polarization properties and can be determined using the data for $N(x_F,H_i|s)$. Since there are no suitable data available for different hyperons, we make the following estimation based on our result for $N_0(x_F,\Lambda|s)$. We simply assume that the x_F dependences of all these $N_0(x_F,H_i|s)$'s are the same. Hence, we have $N_0(x_F,H_i|s)=N_0(x_F,\Lambda|s)\langle n_{H_i}\rangle/\langle n_{\Lambda}\rangle$, where $\langle n_{H_i}\rangle$ is the average number of H_i in the central region. For directly produced hyperons, we take $\langle n_{\Sigma^+}\rangle=\langle n_{\Sigma^-}\rangle=\langle n_{\Sigma^0}\rangle=\langle n_{\Lambda}^{dir}\rangle$, and $\langle n_{\Xi}\rangle=\lambda\langle n_{\Sigma}\rangle$, where $\lambda=0.3$ denotes the strangeness suppression factor. We take Σ^0 and Ξ decays into account, and obtain that $\langle n_{\Lambda}\rangle=(2+2\lambda)\langle n_{\Lambda}^{dir}\rangle$. In this way, we obtain a rough estimation of the $N_0(x_F,H_i|s)$. Using this, we obtain the numerical results for $P_{\Sigma^{\pm}(\text{or }\Xi)}(x_F|s)$ at $p_{inc}=400$ GeV/c as shown in Fig. 2, and those at $p_{inc}=800$ GeV/c in Fig. 3. We see that the results show clearly the qualitative features mentioned above and that these features are in agreement with the data [50–55].

B. Hyperon polarization in $K^-+p\to H_i+X$

There exist also data for hyperon polarization in $K^-p\to\Lambda X$. The data show that P_{Λ} in this process is also significantly different from zero for large x_F. Furthermore, compared with those for P_{Λ} in $pp\to\Lambda X$, the data show that P_{Λ} in $K^-p\to\Lambda X$ has a different sign from that for $pp\to\Lambda X$. For x_F from 0 to 1, P_{Λ} in $K^-p\to\Lambda X$ begins with $P_{\Lambda}\approx 0$ and increases monotonically to about 40% at $x_F\to 1$. Now we apply the proposed picture to this process, compare the results with the data, and make predictions for other hyperons.

1. $K^-+p\to\Lambda+X$

Since we have a meson as projectile, there is no contributing process from group A to $K^-p\to\Lambda X$. There is one contributing process from B, i.e., $s_v^P+(u_sd_s)^T\to\Lambda$. Thus, we have

$$D^{B,s}(x_F,\Lambda|s)=\kappa s_v^K(x^P)f_D(x^T|u_sd_s), \tag{43}$$

where $x^P\approx x_F$ and $x^T\approx m_{\Lambda}^2/(sx_F)$. The superscript K for the quark distribution functions denotes that they are for quarks in a K meson.

Just as $pp\to\Lambda X$, there are also contributions from hyperon decays to $K^-p\to\Lambda X$. In K^-p collisions, we have similar contributing processes from group B to Σ^0 and Ξ production, i.e., $s_v^P+(u_sd_s)^T\to\Sigma^0$, $s_v^P+(u_ss_s)^T\to\Xi^0$ and $s_v^P+(d_ss_s)^T\to\Xi^-$, respectively. The corresponding number densities are given by

$$D^{B,s}(x_F,\Sigma^0|s)=\kappa s_v^K(x^P)f_D(x^T|u_sd_s), \tag{44}$$

$$D^{B,s}(x_F,\Xi^0|s)=\kappa s_v^K(x^P)f_D(x^T|u_ss_s), \tag{45}$$

$$D^{B,s}(x_F,\Xi^-|s)=\kappa s_v^K(x^P)f_D(x^T|d_ss_s), \tag{46}$$

respectively.

Hence, we obtain

$$P_\Lambda(x_F|s)=C\frac{t^F_{\Lambda,s}D^{B,s}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}t^D_{\Lambda,H_j}t^F_{H_j,s}D^{C,H_j}(x_F,\Lambda|s)}{N_0(x_F,\Lambda|s)+D^{B,s}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}D^{C,H_j}(x_F,\Lambda|s)},\tag{47}$$

$$D^{C,H_j}(x_F,\Lambda|s)\approx\frac{m_{H_j}}{E_{\Lambda,H_j}}D^{B,s}\left(\frac{m_{H_j}x_F}{E_{\Lambda,H_j}},H_j|s\right),\tag{48}$$

where the fragmentation spin transfer factors $t^F_{H_j,s}$ for quark flavors s and different hyperons H_j are given in Table I, and the decay spin transfer factors t^D_{Λ,H_j} are given in Sec. III.

Using Eq. (47), we can now calculate $P_\Lambda(x_F|s)$ in K^-p collisions as a function of x_F. Just as mentioned at the end of Sec. III A, for $K^-p\to\Lambda X$, we expect that N_0 dominates at small x_F, while $D^B(x_F,\Lambda|s)$ dominates at large x_F since there is no contributing process from group A. It is clear that $P_\Lambda(x_F|s)>0$ and, if the decay contribution is neglected, $P_\Lambda(x_F|s)=0$ for x_F near 0 and increases to $Ct^F_{\Lambda,s}=C=0.6$ with increasing x_F. The decay contribution should make $P_\Lambda(x_F|s)$ slightly less than C at $x_F\to1$ because $0<t^D_{\Lambda,H_j}t^F_{H_j,s}<t^F_{\Lambda,s}$ for all $H_j=\Sigma^0,\Xi^0,\Xi^-$.

We now use Eq. (47) to obtain the numerical results for $P_\Lambda(x_F|s)$ in K^-p collisions in order to get a more precise feeling for the above-mentioned qualitative features. Here, for the unpolarized quark distribution function of the kaon, we use the GRV-P LO set [56] for that of the pion instead. For $N_0(x_F,H_i|s)$ in this process, we simply assume it to be the same as that in *pp* collisions. The results obtained for $P_\Lambda(x_F|s)$ are given in Fig. 4. We see clearly that, as x_F increases from 0 to 1, the $P_\Lambda(x_F|s)$ obtained starts from 0 and increases to about 40% at $x_F\to1$. These qualitative features are in good agreement with the data [57–61].

2. $K^-+p\to\Sigma^\pm$ (or $\Xi^{0,-}$)$+X$

Similar to $K^-p\to\Lambda X$ in K^-p collisions, there is also only one contributing process from group B to the production of each of these hyperons. For Σ^+ and Σ^-, they are $s_v^P+(u_su_s)^T\to\Sigma^+$ and $s_v^P+(d_sd_s)^T\to\Sigma^-$, respectively. The corresponding number densities are given by

$$D^{B,s}(x_F,\Sigma^+|s)=\kappa s_v^K(x^P)f_D(x^T|u_su_s),\tag{49}$$

$$D^{B,s}(x_F,\Sigma^-|s)=\kappa s_v^K(x^P)f_D(x^T|d_sd_s),\tag{50}$$

respectively. The contributing processes from B to Ξ production have been given in the last subsection and their corresponding number densities are given by Eqs. (45) and (46). Thus, we have

$$P_{\Sigma^+}(x_F|s)=C\frac{t^F_{\Sigma^+,s}D^{B,s}(x_F,\Sigma^+|s)}{N_0(x_F,\Sigma^+|s)+D^{B,s}(x_F,\Sigma^+|s)},\tag{51}$$

$$P_{\Sigma^-}(x_F|s)=C\frac{t^F_{\Sigma^-,s}D^{B,s}(x_F,\Sigma^-|s)}{N_0(x_F,\Sigma^-|s)+D^{B,s}(x_F,\Sigma^-|s)},\tag{52}$$

$$P_{\Xi^0}(x_F|s)=C\frac{t^F_{\Xi^0,s}D^{B,s}(x_F,\Xi^0|s)}{N_0(x_F,\Xi^0|s)+D^{B,s}(x_F,\Xi^0|s)},\tag{53}$$

$$P_{\Xi^-}(x_F|s)=C\frac{t^F_{\Xi^-,s}D^{B,s}(x_F,\Xi^-|s)}{N_0(x_F,\Xi^-|s)+D^{B,s}(x_F,\Xi^-|s)},\tag{54}$$

where the $t^F_{H_i,s}$'s are given in Table I.

It is also clear that, for the productions of these hyperons, N_0 dominates at small x_F, while D^B plays the dominating role at large x_F. We expect that, for $\Sigma^\pm$ in K^-p collisions, $P_{\Sigma^\pm}(x_F|s)=0$ for x_F near 0 and decreases to $Ct^F_{\Sigma,s}=-0.2$ with increasing x_F. For $K^-p\to\Xi X$, $P_\Xi(x_F|s)=0$ for x_F near 0 and increases to $Ct^F_{\Xi,s}=0.4$ with increasing x_F. This means that $P_{\Sigma^\pm}(x_F|s)$ is negative in sign and decreases to

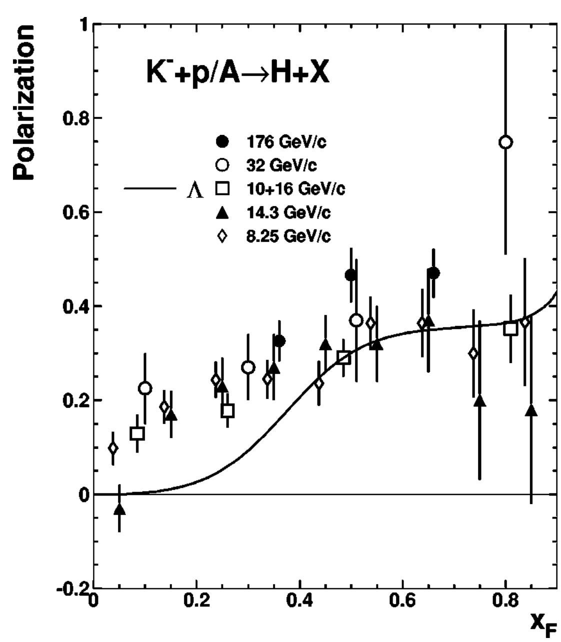

FIG. 4. Calculated results for the polarization of Λ in K^-p collisions as a function of x_F. Data are taken from Refs. [57–61].

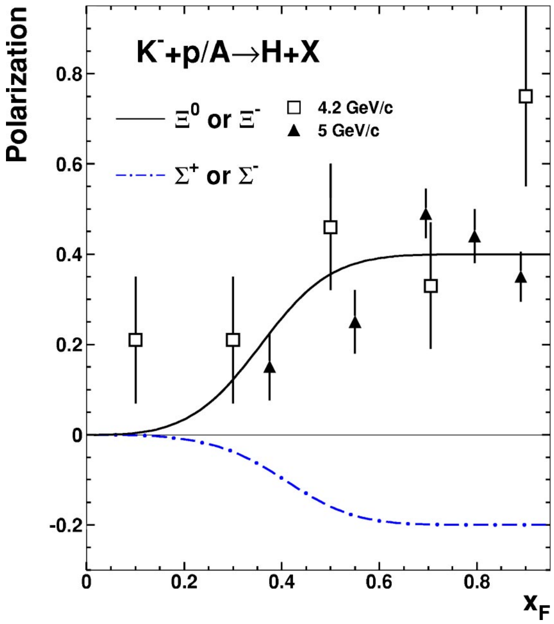

FIG. 5. Calculated results for the polarizations of $\Sigma^{\pm}$ and $\Xi^{0,-}$ in $K^{-}p$ collisions as functions of x_F. Data for Ξ^{-} are taken from Refs. [62,63].

-20% with increasing x_F, while $P_{\Xi}(x_F|s)$ is positive in sign and increases fast to 40% with increasing x_F. The magnitudes of $P_{\Sigma^{\pm}}$ should be smaller than those of P_{Ξ}. These qualitative features are different from those for *pp* collisions and can be checked by future experiments.

By using Eqs. (51)–(54), we also obtain numerical results for $P_{\Sigma^{\pm}}(x_F|s)$ and $P_{\Xi}(x_F|s)$ as functions of x_F in $K^{-}p$ collisions. They are given in Fig. 5. At present, there are data available only for Ξ^{-} [62,63] among these hyperons. We see that the results show clearly the qualitative features mentioned above and that those for Ξ^{-} are consistent with the available data. They can all be further checked by further experiments.

C. Hyperon polarization in $\pi^{\pm}+p\to H_i+X$

Now, we apply the proposed picture to $\pi^{\pm}p$ collisions. From the isospin symmetry, we obtain that P_{Λ} in $\pi^{+}p\to\Lambda X$ is the same as that in $\pi^{-}p\to\Lambda X$, and $P_{\Sigma^{+}}$ and $P_{\Xi^{0}}$ in $\pi^{+}p$ collisions are the same as $P_{\Sigma^{-}}$ and $P_{\Xi^{-}}$ in $\pi^{-}p$ collisions, respectively. So we give only the calculations for $\pi^{-}+p\to H_i+X$.

1. $\pi^{-}+p\to\Lambda+X$

Similar to $K^{-}p\to\Lambda X$, there is also no contributing process from group A to $\pi^{-}p\to\Lambda X$. There is one contributing process from B, i.e., $d_v^P+(u_s s_s)^T\to\Lambda$. Thus, we have

$$D^{B,d}(x_F,\Lambda|s)=\kappa d_v^{\pi}(x^P)f_D(x^T|u_s s_s), \tag{55}$$

where $x^P\approx x_F$ and $x^T\approx m_{\Lambda}^2/(s x_F)$.

Just as in $pp\to\Lambda X$, there are also contributions from hyperon decays to $\pi^{-}p\to\Lambda X$. In $\pi^{-}p$ collisions, we have similar contributing processes from group B to Σ^0 and Ξ^{-} production, i.e., $d_v^P+(u_s s_s)^T\to\Sigma^0$ and $d_v^P+(s_s s_s)^T\to\Xi^{-}$, respectively. Their corresponding number densities are given by

$$D^{B,d}(x_F,\Sigma^0|s)=\kappa d_v^{\pi}(x^P)f_D(x^T|u_s s_s), \tag{56}$$

$$D^{B,d}(x_F,\Xi^{-}|s)=\kappa d_v^{\pi}(x^P)f_D(x^T|s_s s_s), \tag{57}$$

respectively.

Finally, we have

$$P_{\Lambda}(x_F|s)=\frac{C\sum_{H_j=\Sigma^0,\Xi^-}t^D_{\Lambda,H_j}t^F_{H_j,d}D^{C,H_j}(x_F,\Lambda|s)}{N_0(x_F,\Lambda|s)+D^{B,d}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^-}D^{C,H_j}(x_F,\Lambda|s)}, \tag{58}$$

$$D^{C,H_j}(x_F,\Lambda|s)\approx\frac{m_{H_j}}{E_{\Lambda,H_j}}D^{B,d}\left(\frac{m_{H_j}x_F}{E_{\Lambda,H_j}},H_j|s\right), \tag{59}$$

where the fragmentation spin transfer factors $t^F_{H_j,d}$ for quark flavor d and different hyperons H_j are given in Table I, and the decay spin transfer factors t^D_{Λ,H_j} are given in Sec. III.

From Eq. (58), we see immediately that, since $t^F_{\Lambda,d}=0$, $P_{\Lambda}(x_F|s)$ in $\pi p\to\Lambda X$ should be equal to zero if the decay contributions are neglected. The nonzero P_{Λ} in this process comes purely from the decays of heavier hyperons. Taking these decay contributions into account, we expect a small and negative $P_{\Lambda}(x_F|s)$ at large x_F because $Ct^D_{\Lambda,\Sigma^0}t^F_{\Sigma^0,d}=-0.13$ and $Ct^D_{\Lambda,\Xi}t^F_{\Xi^-,d}=-0.18$. The numerical results obtained from Eq. (58) are given in Fig. 6. In the calculations, we use the GRV-P LO set [56] for the unpolarized quark distribution function of pion and simply assume $N_0(x_F,H_i|s)$ in this process to be the same as that in *pp* collisions. We clearly see that, as x_F goes from 0 to 1, the $P_{\Lambda}(x_F|s)$ obtained starts from 0 and decreases to about -10% at large x_F. These qualitative features are in agree-

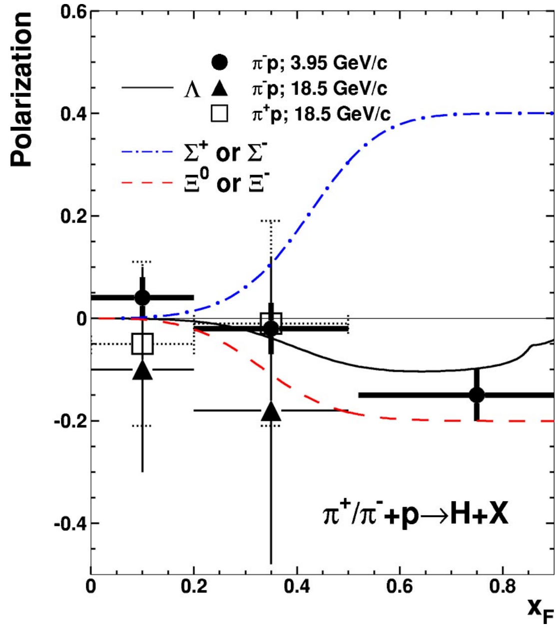

FIG. 6. Calculated results for the polarizations of Λ, $\Sigma^{\pm}$, and $\Xi^{0,-}$ in πp collisions as functions of x_F. Data for Λ are taken from Refs. [64,65].

ment with the data [64,65].

2. $\pi^-+p\rightarrow\Sigma^-$ (or Ξ^-)$+X$

In π^-p collisions, there is one contributing process from group B to the production of Σ^- and Ξ^-. For Σ^-, it is $d_v^P+(d_s s_s)^T\rightarrow\Sigma^-$. The corresponding number density is given by

$$D^{B,d}(x_F,\Sigma^-|s)=\kappa d_v^{\pi}(x^P)f_D(x^T|d_s s_s). \tag{60}$$

The contributing process from B to Ξ^- production was given in the last subsection and its corresponding number density is given by Eq. (57). Thus, we have

$$P_{\Sigma^-}(x_F|s)=C\frac{t^F_{\Sigma^-,d}D^{B,d}(x_F,\Sigma^-|s)}{N_0(x_F,\Sigma^-|s)+D^{B,d}(x_F,\Sigma^-|s)}, \tag{61}$$

$$P_{\Xi^-}(x_F|s)=C\frac{t^F_{\Xi^-,d}D^{B,d}(x_F,\Xi^-|s)}{N_0(x_F,\Xi^-|s)+D^{B,d}(x_F,\Xi^-|s)}. \tag{62}$$

As given in Table I, $t^F_{\Sigma^-,d}=2/3$ and $t^F_{\Xi^-,d}=-1/3$. We thus expect that $P_{\Sigma^-}(x_F|s)=0$ for x_F near 0 and increases to $Ct^F_{\Sigma^-,d}=0.4$ with increasing x_F. For Ξ^-, $P_{\Xi}(x_F|s)=0$ for x_F near 0 and decreases to $Ct^F_{\Xi^-,d}=-0.2$ with increasing x_F. This implies that $P_{\Sigma^-}(x_F|s)$ is positive in sign and increases fast to 40% with increasing x_F, while $P_{\Xi^-}(x_F|s)$ is negative in sign and decreases to -20% with increasing x_F. The magnitude of P_{Σ^-} should be larger than that of P_{Ξ^-}.

Using these equations, we obtain the numerical results for $P_{\Sigma^-}(x_F|s)$ and $P_{\Xi^-}(x_F|s)$ in π^-p collisions as shown in Fig. 6. We see that the results show clearly the qualitative features mentioned above. These features can be tested by future experiments.

D. Hyperon polarization in $\Sigma^-+p\rightarrow H_i+X$

It is interesting to note that experiments on hyperon polarization in reactions using Σ^- beams have also been carried out by the WA89 Collaboration at CERN. Some of the results have already been published [66] and more results are coming. We now apply the proposed picture to this process.

1. $\Sigma^-+p\rightarrow\Lambda+X$

Similar to $pp\rightarrow\Lambda X$, for $\Sigma^-p\rightarrow\Lambda X$, there is one contributing process from group A, i.e., $(d_v s_v)^P+u_s^T\rightarrow\Lambda$, and two contributing processes from B, i.e., $d_v^P+(u_s s_s)^T\rightarrow\Lambda$ and $s_v^P+(u_s d_s)^T\rightarrow\Lambda$. Thus, we have

$$D^A(x_F,\Lambda|s)=\kappa^d_{\Lambda}f_D^{\Sigma^-}(x^P|d_v s_v)u_s(x^T), \tag{63}$$

$$D^{B,d}(x_F,\Lambda|s)=\kappa d_v^{\Sigma^-}(x^P)f_D(x^T|u_s s_s), \tag{64}$$

$$D^{B,s}(x_F,\Lambda|s)=\kappa s_v^{\Sigma^-}(x^P)f_D(x^T|u_s d_s). \tag{65}$$

There are also contributions from hyperon decays to $\Sigma^-p\rightarrow\Lambda X$. For $\Sigma^-p\rightarrow\Sigma^0$(or Ξ^-)X, we have similar contributing processes from groups A and B as those for $\Sigma^-p\rightarrow\Lambda X$. For Σ^0, they are $(d_v s_v)^P+u_s^T\rightarrow\Sigma^0$, $d_v^P+(u_s s_s)^T\rightarrow\Sigma^0$ and $s_v^P+(u_s d_s)^T\rightarrow\Sigma^0$. The number densities of Σ^0 from these processes are given by

$$D^A(x_F,\Sigma^0|s)=\kappa^d_{\Sigma^0}f_D^{\Sigma^-}(x^P|d_v s_v)u_s(x^T), \tag{66}$$

$$D^{B,d}(x_F,\Sigma^0|s)=\kappa d_v^{\Sigma^-}(x^P)f_D(x^T|u_s s_s), \tag{67}$$

$$D^{B,s}(x_F,\Sigma^0|s)=\kappa s_v^{\Sigma^-}(x^P)f_D(x^T|u_s d_s), \tag{68}$$

respectively.

For Ξ^-, they are $(d_v s_v)^P+s_s^T\rightarrow\Xi^-$, $d_v^P+(s_s s_s)^T\rightarrow\Xi^-$, and $s_v^P+(d_s s_s)^T\rightarrow\Xi^-$. The corresponding number densities are given by

$$D^A(x_F,\Xi^-|s)=\kappa^d_{\Xi^-}f_D^{\Sigma^-}(x^P|d_v s_v)s_s(x^T), \tag{69}$$

$$D^{B,d}(x_F,\Xi^-|s)=\kappa d_v^{\Sigma^-}(x^P)f_D(x^T|s_s s_s), \tag{70}$$

$$D^{B,s}(x_F,\Xi^-|s)=\kappa s_v^{\Sigma^-}(x^P)f_D(x^T|d_s s_s), \tag{71}$$

respectively.

For $\Sigma^-p\rightarrow\Xi^0X$, there is one contributing process from B, i.e., $s_v^P+(u_s s_s)^T\rightarrow\Xi^0$. The corresponding number density is given by

$$D^{B,s}(x_F,\Xi^0|s)=\kappa s_v^{\Sigma^-}(x^P)f_D(x^T|u_s s_s). \tag{72}$$

Thus, we obtain that

TABLE IV. Relative weights $w(s_d,s_{dn}|d_v s_v)$ for the $d_v s_v$ diquark from Σ^- in different spin states, those for the produced Λ and Σ^0, the corresponding α_Λ^A and $\alpha_{\Sigma^0}^A$, and the total weights w_Λ^A and $w_{\Sigma^0}^A$.

Possible spin states	$(d_v s_v)_{0,0}u_s^\downarrow$		$(d_v s_v)_{1,0}u_s^\downarrow$		$(d_v s_v)_{1,1}u_s^\downarrow$	
$w(s_d,s_{dn}\|d_v s_v)$	3/4		1/12		1/6	
Possible products	$\Lambda^\downarrow$	$\Sigma^{0\downarrow}$	$\Lambda^\downarrow$	$\Sigma^{0\downarrow}$	$\Lambda^\uparrow$	$\Sigma^{0\uparrow}$
$\|\langle (q_v q_v)_{s_d,s_{dn}} q_s^\downarrow \| H_i(s_n)\rangle\|^2$	1/4	3/4	1/4	1/12	1/2	1/6
The final relative weights	3/16	9/16	1/48	1/144	1/12	1/36
The resulting w_{Hi}^A and α_{Hi}^A	Λ: 7/24, $-3/7$;			Σ^0: 43/72, $-39/43$		

$$N(x_F,\Lambda|s)=N_0(x_F,\Lambda|s)+D^A(x_F,\Lambda|s)+\sum_{f=d,s}D^{B,f}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}D^{C,H_j}(x_F,\Lambda|s), \quad (73)$$

$$D^{C,\Sigma^0}(x_F,\Lambda|s)\approx\frac{m_\Sigma}{E_{\Lambda,\Sigma^0}}\left[D^A\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)+\sum_{f=d,s}D^{B,f}\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)\right], \quad (74)$$

$$D^{C,\Xi^0}(x_F,\Lambda|s)\approx\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,s}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^0|s\right), \quad (75)$$

$$D^{C,\Xi^-}(x_F,\Lambda|s)\approx\frac{m_\Xi}{E_{\Lambda,\Xi}}\left[D^A\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right)+\sum_{f=d,s}D^{B,f}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right)\right]. \quad (76)$$

In the same way as that for $pp\to\Lambda X$, we calculate w_Λ^A, the total relative weight for the production of Λ from group A, and α_Λ^A. The results are shown in Table IV. Similarly, we obtain the corresponding results for Σ^0 production as shown in the same table and those for Ξ^- production as shown in Table V. From these results, we obtain

$$\Delta N(x_F,\Lambda|s)=C\alpha_\Lambda^A D^A(x_F,\Lambda|s)+Ct_{\Lambda,s}^F D^{B,s}(x_F,\Lambda|s)+\sum_{H_j=\Sigma^0,\Xi^0,\Xi^-}\Delta D^{C,H_j}(x_F,\Lambda|s), \quad (77)$$

$$\Delta D^{C,\Sigma^0}(x_F,\Lambda|s)\approx Ct_{\Lambda,\Sigma^0}^D\frac{m_\Sigma}{E_{\Lambda,\Sigma^0}}\left[\alpha_{\Sigma^0}^A D^A\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)+\sum_{f=d,s}t_{\Sigma^0,f}^F D^{B,f}\left(\frac{m_\Sigma x_F}{E_{\Lambda,\Sigma^0}},\Sigma^0|s\right)\right], \quad (78)$$

$$\Delta D^{C,\Xi^0}(x_F,\Lambda|s)\approx Ct_{\Lambda,\Xi}^D t_{\Xi^0,s}^F\frac{m_\Xi}{E_{\Lambda,\Xi}}D^{B,s}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^0|s\right), \quad (79)$$

$$\Delta D^{C,\Xi^-}(x_F,\Lambda|s)\approx Ct_{\Lambda,\Xi}^D\frac{m_\Xi}{E_{\Lambda,\Xi}}\left[\alpha_{\Xi^-}^A D^A\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right)+\sum_{f=d,s}t_{\Xi^-,f}^F D^{B,f}\left(\frac{m_\Xi x_F}{E_{\Lambda,\Xi}},\Xi^-|s\right)\right], \quad (80)$$

where the fragmentation spin transfer factors $t_{H_j,f}^F$ for different quark flavors f and different hyperons H_j are given in Table I, and the decay spin transfer factors t_{Λ,H_j}^D are given in Sec. III.

Using the results given by Eqs. (63)–(80), we can now calculate $P_\Lambda(x_F|s)$ in $\Sigma^- p$ collisions as a function of x_F. Between the two contributing processes from B, the contribution from $s_v^P+(u_s d_s)^T\to\Lambda$ is more important than that from $d_v^P+(u_s s_s)^T\to\Lambda$ due to the strangeness suppression in the sea quarks of the target. We note that $\alpha_\Lambda^A=-3/7$, $t_{\Lambda,d}^F=0$, and $t_{\Lambda,s}^F=1$. Hence, if we neglect the decay contributions, we expect that, for x_F increasing from 0 to 1, $P_\Lambda(x_F|s)$ starts from 0, increases to some positive value, and then decreases to $C\alpha_\Lambda^A=-3C/7$ at $x_F\to1$. We also note that $\alpha_{\Sigma^0}^A=-39/43$, which is large in magnitude and has the same

TABLE V. Relative weights $w(s_d,s_{dn}|d_v s_v)$ for the $d_v s_v$ diquark from Σ^- in different spin states, those for the produced Ξ^-, the corresponding $\alpha_{\Xi^-}^A$, and the total weight $w_{\Xi^-}^A$.

Possible spin states	$(d_v s_v)_{0,0}s_s^\downarrow$	$(d_v s_v)_{1,0}s_s^\downarrow$	$(d_v s_v)_{1,1}s_s^\downarrow$
$w(s_d,s_{dn}\|d_v s_v)$	3/4	1/12	1/6
Possible products	$\Xi^{-\downarrow}$	$\Xi^{-\downarrow}$	$\Xi^{-\uparrow}$
$\|\langle (q_v q_v)_{s_d,s_{dn}} q_s^\downarrow \| H_i(s_n)\rangle\|^2$	3/4	1/12	1/6
The final relative weights	9/16	1/144	1/36
The resulting w_{Hi}^A and α_{Hi}^A		Ξ^-: 43/72, $-39/43$	

sign as α_Λ^A. This is quite different from the situation in $pp \to \Lambda X$ where $\alpha_{\Sigma^0}^A$ has a different sign and smaller magnitude compared to α_Λ^A. Since the decay spin transfer factor $t_{\Lambda,\Sigma^0}^D = -1/3$, we thus expect a very significant contribution from the Σ^0 decay to P_Λ in $\Sigma^- p \to \Lambda X$, and this contribution cancels that from the directly produced Λ at large x_F. As a consequence, taking the decay contributions into account, in particular the significant contribution from $\Sigma^0 \to \Lambda\gamma$, we expect that, for x_F going from 0 to 1, $P_\Lambda(x_F|s)$ starts from 0, increases slowly to some positive value at moderate x_F, and then begins to decrease at some x_F and finally reaches some negative value; but the magnitude is less than $3C/7$ as $x_F \to 1$.

To get some quantitative feeling, we now use Eqs. (63)–(80) to calculate $P_\Lambda(x_F|s)$ in $\Sigma^- p$ collisions numerically. Since our purpose is to get a feeling for the x_F dependence of $P_\Lambda(x_F|s)$, we simply make the following simplifications. For the unpolarized quark distribution functions in Σ^-, we use those in the proton [45] as approximations. $N_0(x_F,H_i|s)$, is taken as the same as that in pp collisions. The numerical results obtained for $P_\Lambda(x_F|s)$ are given in Fig. 7. We see clearly that, as x_F increases from 0 to 1, the $P_\Lambda(x_F|s)$ obtained first increases from 0 to some positive value, and then decreases even to negative at $x_F \to 1$. These qualitative features are just those expected above. They can be checked by further experiments.

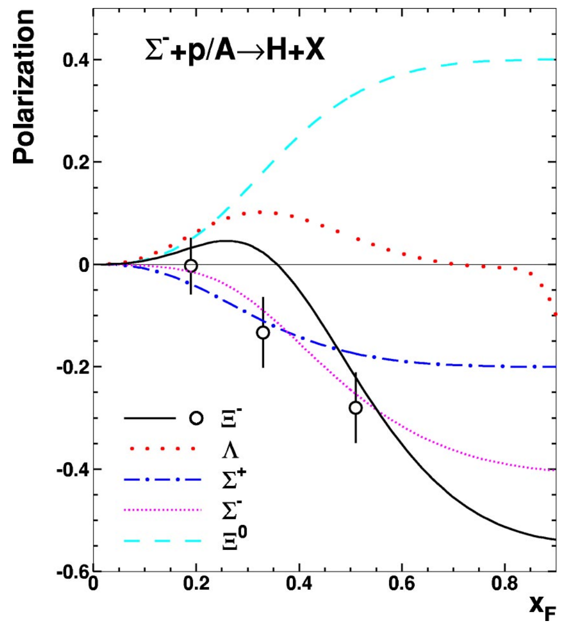

FIG. 7. Calculated results for the polarizations of Λ, $\Sigma^\pm$, and $\Xi^{0,-}$ as functions of x_F. The thin dotted curve represents the polarization of Λ. Data for Ξ^- are taken from Ref. [66].

2. $\Sigma^- + p \to \Sigma^\pm$ (or $\Xi^{0,-}$) $+ X$

First, we look at Σ^- production in $\Sigma^- p$ collisions. For $\Sigma^- p \to \Sigma^- X$, there are two contributing processes from group A, i.e., $(d_v s_v)^P + d_s^T \to \Sigma^-$ and $(d_v d_v)^P + s_s^T \to \Sigma^-$, and two from B, i.e., $d_v^P + (d_s s_s)^T \to \Sigma^-$ and $s_v^P + (d_s d_s)^T \to \Sigma^-$. The corresponding number densities are given by

$$D^{A,ds}(x_F,\Sigma^-|s) = \kappa_{\Sigma^-}^d f_D^{\Sigma^-}(x^P|d_v s_v) d_s(x^T), \tag{81}$$

$$D^{A,dd}(x_F,\Sigma^-|s) = \kappa_{\Sigma^-}^d f_D^{\Sigma^-}(x^P|d_v d_v) s_s(x^T), \tag{82}$$

$$D^{B,d}(x_F,\Sigma^-|s) = \kappa d_v^{\Sigma^-}(x^P) f_D(x^T|d_s s_s), \tag{83}$$

$$D^{B,s}(x_F,\Sigma^-|s) = \kappa s_v^{\Sigma^-}(x^P) f_D(x^T|d_s d_s). \tag{84}$$

Hence, we have

$$P_{\Sigma^-}(x_F|s) = C \frac{\sum_{f=ds,dd} \alpha_{\Sigma^-}^{A,f} D^{A,f}(x_F,\Sigma^-|s) + \sum_{f=d,s} t_{\Sigma^-,f}^F D^{B,f}(x_F,\Sigma^-|s)}{N_0(x_F,\Sigma^-|s) + \sum_{f=ds,dd} D^{A,f}(x_F,\Sigma^-|s) + \sum_{f=d,s} D^{B,f}(x_F,\Sigma^-|s)}, \tag{85}$$

where the $t_{\Sigma^-,f}^F$'s are given in Table I; $\alpha_{\Sigma^-}^{A,f}$ and other related quantities describing Σ^- from the two contributing processes from A are shown in Tables VI and VII, respectively.

Because of the strangeness suppression in the sea quarks of the target, we expect that the most important contributing process to the production of Σ^- from A is $(d_v s_v)^P + d_s^T \to \Sigma^-$ and that from B is $s_v^P + (d_s d_s)^T \to \Sigma^-$. From Tables VI and I, we see that $\alpha_{\Sigma^-}^{A,ds} = -39/43$ and $t_{\Sigma^-,s}^F = -1/3$. Both of them contribute negatively to the polarization of Σ^- in $\Sigma^- p \to \Sigma^- X$. On the contrary, from Tables VII and I, we see that $\alpha_{\Sigma^-}^{A,dd} = 3/5$ and $t_{\Sigma^-,d}^F = 2/3$. Both of them are positive and $\alpha_{\Sigma^-}^{A,dd} < |\alpha_{\Sigma^-}^{A,ds}|$, but $t_{\Sigma^-,d}^F = 2|t_{\Sigma^-,s}^F|$. We thus expect that a large portion of the contribution from $s_v^P + (d_s d_s)^T \to \Sigma^-$ to P_{Σ^-} is canceled by that from $d_v^P + (d_s s_s)^T \to \Sigma^-$, while a relatively small fraction of the contribution from $(d_v s_v)^P + d_s^T \to \Sigma^-$ is canceled by that from $(d_v d_v)^P + s_s^T \to \Sigma^-$. Hence, for x_F from 0 to 1, $P_{\Sigma^-}(x_F|s)$ in $\Sigma^- p \to \Sigma^- X$ should start from 0, decrease slowly to some negative value at moderate x_F, and then tend to a result between $C(\alpha_{\Sigma^-}^{A,ds} + \alpha_{\Sigma^-}^{A,dd}) = -0.18$ and $C\alpha_{\Sigma^-}^{A,ds} = -0.54$ at $x_F \to 1$.

TABLE VI. Relative weights $w(s_d, s_{dn}|d_v s_v)$ for the $d_v s_v$ diquark from Σ^- in different spin states, those for the produced Σ^-, the corresponding $\alpha_{\Sigma^-}^{A,ds}$, and the total weight $w_{\Sigma^-}^{A,ds}$.

Possible spin states	$(d_v s_v)_{0,0} d_s^{\downarrow}$	$(d_v s_v)_{1,0} d_s^{\downarrow}$	$(d_v s_v)_{1,1} d_s^{\downarrow}$
$w(s_d, s_{dn}\|d_v s_v)$	3/4	1/12	1/6
Possible products	$\Sigma^{-\downarrow}$	$\Sigma^{-\downarrow}$	$\Sigma^{-\uparrow}$
$\|\langle (q_v q_v)_{s_d, s_{dn}} q_s^{\downarrow} \| H_i(s_n)\rangle\|^2$	3/4	1/12	1/6
The final relative weights	9/16	1/144	1/36
The resulting w_{Hi}^A and α_{Hi}^A		Σ^-: 43/72, $-39/43$	

Using Eq. (85), we calculate $P_{\Sigma^-}(x_F|s)$ in this process. The numerical results for it are given in Fig. 7. We see that $P_{\Sigma^-}(x_F|s)$ indeed starts from 0, decreases to about -20% at moderate x_F, and reaches about -40% at $x_F \to 1$. These features are the same as those from the qualitative analysis and can be tested by future experiments.

Then we look at the production of Σ^+, Ξ^0, and Ξ^- in $\Sigma^- p$ collisions. For Σ^+, there is no contributing process from group A but one contributing process from B, i.e., $s_v^P + (u_s u_s)^T \to \Sigma^+$. The corresponding number density is given by

$$D^{B,s}(x_F, \Sigma^+|s) = \kappa s_v^{\Sigma^-}(x^P) f_D(x^T|u_s u_s). \tag{86}$$

The contributing processes to Ξ production have been given in the last subsection and their corresponding number densities are given by Eqs. (69)–(72). Hence, we have

$$P_{\Sigma^+}(x_F|s) = C \frac{t_{\Sigma^+,s}^F D^{B,s}(x_F, \Sigma^+|s)}{N_0(x_F, \Sigma^+|s) + D^{B,s}(x_F, \Sigma^+|s)}, \tag{87}$$

$$P_{\Xi^0}(x_F|s) = C \frac{t_{\Xi^0,s}^F D^{B,s}(x_F, \Xi^0|s)}{N_0(x_F, \Xi^0|s) + D^{B,s}(x_F, \Xi^0|s)}, \tag{88}$$

$$P_{\Xi^-}(x_F|s) = C \frac{\alpha_{\Xi^-}^A D^A(x_F, \Xi^-|s) + \sum_{f=d,s} t_{\Xi^-,f}^F D^{B,f}(x_F, \Xi^-|s)}{N_0(x_F, \Xi^-|s) + D^A(x_F, \Xi^-|s) + \sum_{f=d,s} D^{B,f}(x_F, \Xi^-|s)}, \tag{89}$$

where the $t_{H_i,f}^F$'s are given in Table I and $\alpha_{\Xi^-}^A$ is given in Table V.

We recall that $t_{\Sigma^+,s}^F = -1/3$ and $t_{\Xi^0,s}^F = 2/3$. Thus, $P_{\Sigma^+}(x_F|s)$ should start from 0 and decrease to $Ct_{\Sigma^+,s}^F = -0.2$ with increasing x_F, while $P_{\Xi^0}(x_F|s)$ begins from 0 and increases to $Ct_{\Xi^0,s}^F = 0.4$. This means that P_{Σ^+} in this process is negative in sign and relatively small in magnitude but P_{Ξ^0} is positive and large. For Ξ^- production, we have $\alpha_{\Xi^-}^A = -39/43$, $t_{\Xi^-,s}^F = 2/3$, and $t_{\Xi^-,d}^F = -1/3$. Hence, we expect from Eq. (89) that, as x_F increases from 0 to 1, $P_{\Xi^-}(x_F|s)$ starts from 0, increases to some positive value below $Ct_{\Xi^-,s}^F = 0.4$, begins to decrease at some x_F, and fi-

TABLE VII. Relative weights $w(s_d, s_{dn}|d_v d_v)$ for the $d_v d_v$ diquark from Σ^- in different spin states, those for the produced Σ^-, the corresponding $\alpha_{\Sigma^-}^{A,dd}$, and the total weight $w_{\Sigma^-}^{A,dd}$.

Possible spin states	$(d_v d_v)_{0,0} s_s^{\downarrow}$	$(d_v d_v)_{1,0} s_s^{\downarrow}$	$(d_v d_v)_{1,1} s_s^{\downarrow}$
$w(s_d, s_{dn}\|d_v d_v)$	0	1/3	2/3
Possible products	$\Sigma^{-\downarrow}$	$\Sigma^{-\downarrow}$	$\Sigma^{-\uparrow}$
$\|\langle (q_v q_v)_{s_d, s_{dn}} q_s^{\downarrow} \| H_i(s_n)\rangle\|^2$	—	1/3	2/3
The final relative weights	—	1/9	4/9
The resulting w_{Hi}^A and α_{Hi}^A		Σ^-: 5/9, 3/5	

nally reaches $C\alpha^{A}_{\Xi^-}=-0.54$. In Fig. 7, we show the obtained numerical results as functions of x_F. We see clearly that the results indeed show the above-mentioned qualitative features. These features can be checked by future experiments.

We emphasize again that the most important purpose of the numerical results presented in this section is to show the qualitative features of hyperon polarization in different reactions obtained in the proposed picture. For this purpose, we make several simplifications to reduce the free parameters in connection with the number densities in unpolarized reactions. Further improvements can be made if more accurate data are available.

V. SUMMARY AND OUTLOOK

In summary, we have calculated the polarizations of different hyperons as functions of x_F in the inclusive pp, K^-p, $\pi^{\pm}p$, and Σ^-p collisions. We used the picture proposed in a previous Letter [10], which relates the hyperon polarization in unpolarized hadron-hadron collisions to the left-right asymmetry in singly polarized reactions. We discussed the qualitative features for hyperon polarizations in these reactions and presented the corresponding numerical results. These qualitative features are all in agreement with the available data. Predictions for future experiments have been made. These predictions can be used as good tests for the picture.

It should be emphasized that several points need to be further developed in the model. One of them is the influence of the vector meson production associated with the hyperon that contains a valence diquark of the projectile as mentioned at the end of Sec. II. Such a study is under way. Another very important aspect is the transverse momentum dependence of the hyperon polarization. It is clear that the general formulas given in Sec. III can be extended to include the p_T dependence. We can use them to calculate the p_T dependence of P_H in the model. This is also a very important aspect of the existing data and should be taken as a further challenge to the model. As can be seen from the formulas in Sec. III, the p_T dependence of P_H should come from the p_T dependence of C and the interplay of the p_T-dependent N_0 and D. The p_T dependence of the direct-formation contribution D comes mainly from those of the quark distributions. The N_0 part can be parametrized from the data on unpolarized cross sections. But to get the p_T dependence of C we need data on the p_T dependence of A_N. Presently, there are no such data available. We get only an average C in the p_T interval $0.7<p_T<2.0$ GeV/c. This is the largest difficulty in calculating the p_T dependence of P_H at the present stage. Nevertheless, a phenomenological analysis can and should be made. Such studies are under way. The results we obtained in Sec. IV should be taken as the average results in the corresponding p_T region. Since many of the data are from fixed angle experiments, the comparison of our results with the data can only be regarded as qualitative.

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[1] A. Lesnik *et al.*, Phys. Rev. Lett. **35**, 770 (1975).
[2] G. Bunce *et al.*, Phys. Rev. Lett. **36**, 1113 (1976).
[3] For a review of the data, see, e.g., K. Heller, in *Proceedings of the 12th International Symposium on High Energy Spin Physics*, Amsterdam, Netherlands, 1996, edited by C.W. de Jager, T.J. Ketel, P.J. Mulders, J.E.J. Oberski, and M. Oskam-Tamboezer (World Scientific, Singapore, 1997), p. 23, and the references given there.
[4] B. Andersson, G. Gustafson, and G. Ingelman, Phys. Lett. **85B**, 417 (1979).
[5] T.A. DeGrand and H.I. Miettinen, Phys. Rev. D **24**, 2419 (1981).
[6] J. Szwed, Phys. Lett. **105B**, 403 (1981).
[7] L.G. Pondrom, Phys. Rep. **122**, 57 (1985).
[8] R. Barni, G. Preparata, and P.G. Ratcliffe, Phys. Lett. B **296**, 251 (1992).
[9] J. Soffer and N.A. Törnqvist, Phys. Rev. Lett. **68**, 907 (1992).
[10] Liang Zuo-tang and C. Boros, Phys. Rev. Lett. **79**, 3608 (1997).
[11] Liang Zuo-tang and C. Boros, Phys. Rev. D **61**, 117503 (2000).
[12] Liang Zuo-tang and Li Tie-shi, J. Phys. G **26**, L111 (2000).
[13] M. Anselmino, D. Boer, U. D'Alesio, and F. Murgia, Phys. Rev. D **63**, 054029 (2001).
[14] A summary of the data can be found in, e.g., A. Bravar, in *Proceedings of the 13th International Symposium on High Energy Spin Physics*, Protvino, Russia, 1998, edited by N.E. Tyurin *et al*. (World Scientific, Singapore, 1999), p. 167; see in particular also Ref. [15] listed below.
[15] FNAL E704 Collaboration, D.L. Adams *et al.*, Phys. Lett. B **264**, 462 (1991); **276**, 531 (1992); Z. Phys. C **56**, 181 (1992); A. Bravar *et al.*, Phys. Rev. Lett. **75**, 3073 (1995); **77**, 2626 (1996); D.L. Adams *et al.*, Nucl. Phys. **B510**, 3 (1998).
[16] Liang Zuo-tang and Xu Qing-hua, Phys. Rev. D **65**, 113012 (2002).
[17] R608 Collaboration, T. Henkes, Phys. Lett. B **283**, 155 (1992).
[18] CLAS Collaboration, D.S. Carman *et al.*, Phys. Rev. Lett. **90**, 131804 (2003).
[19] C. Boros and Liang Zuo-tang, Phys. Rev. D **57**, 4491 (1998).
[20] G. Gustafson and J. Häkkinen, Phys. Lett. B **303**, 350 (1993).
[21] ALEPH Collaboration, D. Buskulic *et al.*, Phys. Lett. B **374**, 319 (1996).
[22] OPAL Collaboration, K. Ackerstaff *et al.*, Eur. Phys. J. C **2**, 49 (1998).
[23] Liu Chun-xiu and Liang Zuo-tang, Phys. Rev. D **62**, 094001 (2000); Liu Chun-xiu, Xu Qing-hua, and Liang Zuo-tang, *ibid.* **64**, 073004 (2001); Xu Qing-hua, Liu Chun-xiu, and Liang

Zuo-tang, *ibid.* **65**, 114008 (2002); Liang Zuo-tang and Liu Chun-xiu, *ibid.* **66**, 057302 (2002).
[24] NOMAD Collaboration, P. Astier *et al.*, Nucl. Phys. **B588**, 3 (2000).
[25] Liang Zuo-tang and C. Boros, Int. J. Mod. Phys. A **15**, 927 (2000).
[26] C. Boros, Liang Zuo-tang, and Meng Ta-chung, Phys. Rev. Lett. **70**, 1751 (1993); Liang Zuo-tang and Meng Ta-chung, Phys. Rev. D **49**, 3759 (1994).
[27] Xu Qing-hua and Liang Zuo-tang, Phys. Rev. D **68**, 034023 (2003).
[28] P. Chliapnikov *et al.*, Nucl. Phys. **B37**, 336 (1972); I.V. Ajinenko *et al.*, Z. Phys. C **5**, 177 (1980); K. Paler *et al.*, Nucl. Phys. **B96**, 1 (1975); Yu. Arestov *et al.*, Z. Phys. C **6**, 101 (1980).
[29] M. Barth *et al.*, Nucl. Phys. **B223**, 296 (1983).
[30] EXCHARM Collaboration, A.N. Aleev *et al.*, Phys. Lett. B **485**, 334 (2000).
[31] ALEPH Collaboration, D. Buskulic *et al.*, Z. Phys. C **69**, 393 (1995).
[32] DELPHI Collaboration, P. Abreu *et al.*, Z. Phys. C **68**, 353 (1995); Phys. Lett. B **406**, 271 (1997).
[33] OPAL Collaboration, K. Ackerstaff *et al.*, Phys. Lett. B **412**, 210 (1997); Z. Phys. C **74**, 437 (1997); G. Abbiendi *et al.*, Eur. Phys. J. C **16**, 61 (2000).
[34] Xu Qing-hua, Liu Chun-xiu, and Liang Zuo-tang, Phys. Rev. D **63**, 111301 (2001).
[35] C. Boros and Liang Zuo-tang, Phys. Rev. D **53**, R2279 (1996).
[36] Dong Hui and Liang Zuo-tang (in preparation).
[37] W. Ochs, Nucl. Phys. **B118**, 397 (1977).
[38] W. Kittel, in *Partons in Soft-Hadronic Processes*, Proceedings of the Europhysics Study Conference, Erice, Italy, 1981, edited by R.T. Van de Walle (World Scientific, Singapore, 1981), p. 1.
[39] K.P. Das and R.C. Hwa, Phys. Lett. **68B**, 459 (1977).
[40] Dong Hui, Li Fang-zhen, and Liang Zuo-tang, Phys. Rev. D **69**, 017501 (2004).
[41] E925 Collaboration, C.E. Allgower *et al.*, Phys. Rev. D **65**, 092008 (2002).
[42] R. Gatto, Phys. Rev. **109**, 610 (1958).
[43] Particle Data Group, K. Hagiwara, Phys. Rev. D **66**, 010001 (2002).
[44] In practice, since there are no data directly on $N(x_F,\Lambda|s)$, we determine the two κ's by fitting a few data points at large x_F of the available data on the invariant cross section $Ed^3\sigma/dp^3$ at a given p_T because apparently $Ed^3\sigma/dp^3 \propto x_F N(x_F,s)$. This implies that we multiply the number densities $N(x_F,\Lambda|s)$, $N_0(x_F,\Lambda|s)$, and $D(x_F,\Lambda|s)$'s by x_F and a constant which changes the differential cross sections to the number densities. In this way, all of them are changed to their counterparts in the invariant cross section $Ed^3\sigma/dp^3$. It is expected that, if the direct-fusion model works, κ should strongly depend on neither x nor p_T. The x_F and p_T dependences of the number density of the hadron produced from this process come predominantly from the quarks and/or antiquarks in the initial states. So we can approximately fix the κ parameters entering a p_T integrated density from data at fixed p_T. For $p_T = 0.65$ GeV/c, the corresponding result for $x_F N_0$ is $0.2(1-x_F)^6 e^{-3x_F^5}$ mb/GeV^2, that for κ^d is 1.157×10^{-2} mb/GeV^2, and that for κ is 1.62×10^{-4} mb/GeV^2, For more details, see Refs. [25,26,35] given above.
[45] See, e.g., M. Glück, E. Reya, and A. Vogt, Eur. Phys. J. C **5**, 461 (1998).
[46] See, e.g., J.J. Dugne and P. Tavernier, Z. Phys. C **59**, 333 (1993).
[47] A.M. Smith *et al.*, Phys. Lett. B **185**, 209 (1987).
[48] B. Lundberg *et al.*, Phys. Rev. D **40**, 3557 (1989).
[49] E.J. Ramberg *et al.*, Phys. Lett. B **338**, 403 (1994).
[50] C. Wilkinson *et al.*, Phys. Rev. Lett. **58**, 855 (1987).
[51] L. Deck *et al.*, Phys. Rev. D **28**, 1 (1983).
[52] R. Rameika *et al.*, Phys. Rev. D **33**, 3172 (1986).
[53] K. Heller *et al.*, Phys. Rev. Lett. **51**, 2025 (1983).
[54] A. Morelos *et al.*, Phys. Rev. D **52**, 3777 (1995).
[55] J. Duryea *et al.*, Phys. Rev. Lett. **67**, 1193 (1991).
[56] See, e.g., M. Glück, E. Reya, and A. Vogt, Z. Phys. C **53**, 651 (1992).
[57] S.A. Gourlay *et al.*, Phys. Rev. Lett. **56**, 2244 (1986).
[58] M.L. Faccini-Turluer *et al.*, Z. Phys. C **1**, 19 (1979).
[59] H. Grassler *et al.*, Nucl. Phys. **B136**, 386 (1978).
[60] H. Abramowicz *et al.*, Nucl. Phys. **B105**, 222 (1976).
[61] M. Baubillier *et al.*, Nucl. Phys. **B148**, 18 (1979).
[62] S.N. Ganguli *et al.*, Nucl. Phys. **B128**, 408 (1977).
[63] J. Bensinger *et al.*, Nucl. Phys. **B252**, 561 (1985).
[64] B. Adeva *et al.*, Z. Phys. C **26**, 359 (1984).
[65] D.H. Stuntebeck *et al.*, Phys. Rev. D **9**, 608 (1974).
[66] WA89 Collaboration, M.I. Adamovich *et al.*, Z. Phys. A **350**, 379 (1995).