NN scattering amplitudes from 90° c.m. into the Landshoff region

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It is likely that nucleon-nucleon elastic scattering at large, fixed $|t|$ can be described by the Landshoff three-gluon exchange mechanism. However, a phenomenological normalization rules out the possibility that the Landshoff mechanism is involved in producing the sharp structure of the two-spin asymmetry A_{NN}, observed experimentally at $p_{\rm lab} \approx 13$ GeV/c, or that it is involved in producing the oscillations observed in $d\sigma/dt$ at 90° in the c.m. frame. A simplified analysis of the 90° amplitudes shows that it is likely that some subasymptotic mechanism interfering with the dominant quark-interchange amplitudes is responsible for both these intriguing phenomena. While the large-angle data show evidence for amplitudes approaching those of exclusive QCD, the Landshoff region in $pp \to pp$ shows no evidence for Sudhakov suppression nor evidence for effects associated with the running of the QCD coupling. We suggest that a measurement of the energy dependence of the polarization asymmetry at large $|-t|$ can greatly enlarge the understanding of exclusive QCD.

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I. INTRODUCTION

The study of exclusive hard scattering processes offers a formidable array of theoretical challenges. Existing data on fixed-angle differential cross sections at high energies [1] confirm that the regularities

$$\lim_{s\to\infty} \frac{d\sigma}{dt}(AB \to CD) \approx f(\theta) s^{2-(n_A+n_B+n_C+n_D)}, \quad (1.1)$$

of the constituent scaling rules [2,3] are observed in all measured processes. The success of the constituent counting rules of scale-invariant field theories suggests that the existing data on exclusive hard scattering may be in a kinematic regime where hadronic amplitudes have the potential to be calculated perturbatively. There are enormous technical difficulties in these calculations, but it is important to persist in the attempt to understand the dynamical foundation for exclusive processes in QCD.

A brief menu of some of the problems in the calculation of the amplitudes for $pp \to pp$ scattering can be instructive. The complexity begins with calculation of the Born amplitudes. Because of combinatorial factors, there are more than 300 000 distinguishable Feynman diagrams for $6q \to 6q$ processes, so that enumeration and evaluation of the graphs requires considerable computational power [4,5]. The predictive capability of the leading-order calculation is impaired by the high degree of renormalization group dependence associated with the factor (α_s/π) [5] which appears in the Born term. In fact, since the perturbative expansion is an asymptotic series, there may be intrinsic barriers to a sensible interpretation of higher-order corrections to these specialized multiparticle processes. There also exists a closely related issue concerning the sensitivity of the calculation to the factorization scheme which absorbs the infrared and collinear singularities of the $6q \to 6q$ amplitude into the proton wave functions. The specific definition of the wave function affects the normalization of pinch and end point singularities in ways which are not yet fully understood. The result of this ambiguity is that small changes in the hadronic wave function can be seen to result in large changes in the normalization of an individual scattering amplitude. With the current level of understanding, there is no control over the normalization of subasymptotic (nonleading twist) corrections to the asymptotic behavior of the amplitudes. With all of these problems, there is little chance of a reliable "first principles" calculation for $pp \to pp$ fixed-angle scattering based on the present technology of QCD perturbation theory.

In many respects, the proton-proton process is too complicated for a direct confrontation with this barrage of theoretical problems. Similar scattering processes must be subjected to quantitative calculations in order to test theoretical techniques. For example, detailed analysis of perturbative calculations of the nucleon form factor can shed light on many of the theoretical problems mentioned above. In fact, if the proton form factor at large momentum transfer were well understood theoretically, it is possible that a "reduced-amplitude" formulation based on sophisticated combinatoric tools could provide a consistent normalization for $pp \to pp$ amplitudes.

However, in the absence of a well-defined perturbative calculation, there are valid reasons to continue the study of $pp \to pp$ elastic scattering. High intensity proton beams and a clean signature for the elastic scattering process make this one of the best studied set of exper-

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imental cross sections. Spin-dependent observables and the spin-averaged cross section are accurately measured over a wide range of kinematic variables. In addition, most of the normalization problems discussed above are not sensitive to quark combinations. This implies that the ratio of helicity-conserving amplitudes can be predicted [6]. For a given Jacob-Wick helicity amplitude [7], it is possible to do a phenomenological analysis where the normalization of the amplitude is fit to data at one point while its s and t dependence are extracted from the theoretical calculation. This approach leads to a large number of asymptotic predictions and has not yet been fully exploited.

The reason for this omission may be that there has been some confusion over which mechanisms are important. In Sec. II we review how the normalization of the Landshoff three-gluon mechanism amplitude [8] can be reliably specified at high energy and fixed-t from elastic scattering data at the CERN Intersecting Storage Rings (ISR) and the CERN collider. From the known kinematic dependence of the Landshoff amplitude, this phenomenological normalization definitely rules out the possibility that the Landshoff mechanism is involved in either the sharp structure of the A_{NN} data in large angle scattering or the oscillations of the spin-averaged cross section observed at 90° in the c.m. system [9–12]. A consistent normalization of the theoretical amplitudes demands that the scattering is dominated by a "quark interchange" mechanism in kinematic regions where sharp structure has been observed. In Sec. III, we show that if we assume the basic validity of a quark-interchange model (QIM) approach at large angles, it is possible to use the symmetries which occur at 90° in the c.m. system to do a simple amplitude analysis which can extract the basic features of the subasymptotic mechanism responsible for the structure observed. This may have important ramifications for studies of nuclear transparency and for the study of other exclusive processes.

We argue that the structure of the elastic amplitudes within the framework of the constituent-based hard-scattering model can be severely constrained by measurements of the elastic polarization asymmetry at large s and for $|t| \geq 4$ GeV^2. In the traditional Regge theory approach to asymptotic amplitudes, the factorization of Regge residues means that the polarization asymmetry vanishes at large s. In contrast, the hard-scattering approach allows for possible helicity-flip effects associated with the hadronic wave function. These effects vanish at large $|-t|$, but the helicity-flip amplitudes can share the same s dependence as the helicity-conserving ones. This leads to a polarization asymmetry which is almost s independent and hence, falls off only as a function of t [13].

Looking in more detail at the shape of the cross section for $pp \to pp$ in the Landshoff region, it is interesting to note that the data give no support to a running QCD coupling. There is also no evidence for t-dependent suppression associated with a Sudhakov form factor [14]. The search for these effects in the data must be considered one of the challenges for the construction of a reliable phenomenology for exclusive hadronic processes.

II. PHENOMENOLOGICAL NORMALIZATION OF LANDSHOFF AMPLITUDES

We shall define observables for $NN \to NN$ elastic scattering in terms of the following Jacob-Wick helicity amplitudes:

$$\begin{aligned} \Phi_1(s,t) &= \langle ++|M|++\rangle, \\ \Phi_2(s,t) &= \langle ++|M|--\rangle, \\ \Phi_3(s,t) &= \langle +-|M|+-\rangle, \\ \Phi_4(s,t) &= \langle +-|M|-+\rangle, \\ \Phi_5(s,t) &= \langle ++|M|+-\rangle. \end{aligned} \tag{2.1}$$

Other helicity amplitudes are related to these five independent amplitudes using parity conservation, time-reversal invariance, and identical particle symmetry.

The differential cross section is given in terms of these amplitudes as

$$\begin{aligned} \frac{d\sigma}{dt} = \frac{\pi}{2s(s-4m^2)}[&|\Phi_1|^2 + |\Phi_2|^2 + |\Phi_3|^2 \\ &+|\Phi_4|^2 + 4|\Phi_5|^2], \end{aligned} \tag{2.2}$$

while the expressions for the spin dependent observables are given in terms of these amplitudes by [13]

$$\begin{aligned} \Sigma &= s(s-4m^2)\frac{d\sigma}{dt}, \\ P\Sigma &= -\text{Im}[\Phi_5^*(\Phi_1+\Phi_2+\Phi_3-\Phi_4)], \\ A_{SL}\Sigma &= -\text{Re}[\Phi_5^*(\Phi_1+\Phi_2-\Phi_3+\Phi_4)], \\ A_{NN}\Sigma &= \text{Re}[\Phi_1\Phi_2^* - \Phi_3\Phi_4^* + 2|\Phi_5|^2], \\ A_{SS}\Sigma &= \text{Re}[\Phi_1\Phi_2^* + \Phi_3\Phi_4^*], \\ A_{LL}\Sigma &= \frac{1}{2}[|\Phi_3|^2 + |\Phi_4|^2 - |\Phi_1|^2 - |\Phi_2|^2]. \end{aligned} \tag{2.3}$$

In discussing these observables, we make the basic assumption that we can separate a soft, coherent Regge contribution, $\Phi_i^R(s,t)$, for each independent amplitude. These Regge amplitudes dominate the observables at small t. We also assume there exists a "hard" component for each amplitude which obeys the Brodsky-Lepage factorization [15] at large t. For the hard constituent-based component, it is convenient to separate the amplitudes by the number of quarks exchanged in the t channel. For example, the Landshoff mechanism corresponds to three-gluon exchange with no quarks in the t channel [8]. By crossing symmetry, we also have to include three-gluon exchange in the u channel and hence, three-quarks and three-antiquarks in the t channel. The quark interchange mechanism corresponds to either $q\bar{q}$ or $q\bar{q}q\bar{q}$ in the t channel. For each helicity amplitude, we therefore have the decomposition

$$\Phi_i(s,t) = \Phi_i^R(s,t) + \Phi_i^L(s,t) + \Phi_i^Q(s,t) \quad (i=1,\ldots,5), \tag{2.4}$$

where R, L, and Q stand for "Regge," "Landshoff," and "quark interchange," respectively. It is important to note that the separation suggested in (2.4) is largely a matter of convention. The coherent Regge components should be exponentially suppressed at large t, reflecting the size of the individual proton. However, it is possible to absorb part of the coherent portion of the amplitude into the hard, constituent-based component and change the extrapolation of the hard component to finite t values. This is one of the necessary complications in the process of combining formulas applicable in different asymptotic regions, which preclude an unambiguous theoretical normalization of the various components. With our present knowledge, it is necessary to fit the normalization to data. The s and t dependence of the Landshoff and QIM amplitudes are specified by simple QCD calculations, and the assumption we will use for our calculations is that the proton is well described by its minimal three-quark Fock state. For $pp \to pp$, the Landshoff amplitudes at large t and small angle ($\frac{-t}{s} \to 0$) can be written in the form

$$\Phi_1^L \approx \frac{sL(t)}{(2m_p^2-t)^4}[1+O(t/s)],$$
$$\Phi_3^L \approx \frac{sL(t)}{(2m_p^2-t)^4}[1+O(t/s)]. \tag{2.5}$$

Since Φ_4^L is dominated by the u-channel contributions we assume that $\Phi_4^L \approx 0$ in this region. For convenience, we also parametrize the helicity-flip amplitudes as

$$\Phi_5^L \approx -\epsilon(-t)^{\frac{1}{2}} \frac{2sL(t)}{(2m_p^2-t)^5}[1+O(t/s)], \tag{2.6}$$

with the double spin-flip amplitude, $\phi_2^L \approx 0$, since it is down by an additional factor of $(-t)^{-1/2}$ at large $-t$. In Eqs. (2.5) and (2.6), $L(t)$ is a normalization factor which involves a convolution over the proton wave function. The parameter ϵ in (2.6) is assumed to be small, reflecting the size of the known polarization data [14]. The physical significance of this parameter will be discussed in Sec. IV. In the absence of effects associated with Sudhakov form factors [15], $L(t)$ would be constant at large $|t|$. However, it is expected that asymptotically, $L(t)$ shows the behavior [12]

$$L(t) \approx ct^{-\beta}, \tag{2.7}$$

where β is a slowly varying function of t and $-t/s$. We will not assume a specific form for $L(t)$ in this work. Combining (2.2) and (2.5) in a region where the Landshoff amplitudes dominate, (i.e., $m_p^2 \ll -t \ll s$) we have the phenomenological expression

$$\frac{d\sigma}{dt} = \pi \frac{|L(t)|^2}{(2m_p^2-t)^8}\left[2+8|\epsilon|^2\frac{(-t)}{(2m_p^2-t)^2}+\cdots\right], \tag{2.8}$$

where the second term in square brackets is associated with $|\Phi_5|^2$ and other terms are suppressed by additional powers of $\frac{t}{s}$.

Comparing (2.8) with data at large s and with $-t$ outside the coherent Regge region is the first step in a phenomenological normalization. Donnachie and Landshoff [17] have done a thorough phenomenological study of the differential cross sections for pp and $p\bar{p}$ at high energies. We shall review here some aspects of that analysis in order to fix our notation. In the approach of Ref. [17], the dip in the differential cross section at $|-t| = 1.3$ GeV^2 is understood as an interference between double Pomeron exchange and the triple-gluon exchange amplitudes given in (2.5). At larger $-t$, ($|-t| \geq 4$ GeV^2), the three-gluon exchange mechanism dominates. In this kinematic regime, there is strong experimental support for an approximately energy-independent component of the cross section which behaves as t^{-8}. This can be seen in Fig. 1 where we plot $t^8(d\sigma/dt)$ for different energies. Note that the t^{-8} behavior persists to the largest values measured, so that there is a region where (2.8) is valid. Experimentally, the numerical value extracted from these data is

$$|L(t)|^2\Big|_{-t\approx 5\ \text{GeV}^2} \cong 30\,\text{mb GeV [14]}, \tag{2.9}$$

which can be used to normalize this mechanism.

A significant confirmation of the identification of the triple-gluon Landshoff mechanism in the data comes from a comparison of the pp and $p\bar{p}$ channels. The Pomeron and two-Pomeron contributions to the amplitudes have even charge conjugation while the three-gluon mechanism has odd charge conjugation. The interference effect leading to the dip structure in the pp data should therefore not be present in the $p\bar{p}$ channel, since the Landshoff amplitude will change sign. The $p\bar{p}$ measurements at the CERN Super Proton Synchrotron ($Sp\bar{p}S$) collider have confirmed these expectations. No dip is seen in $p\bar{p}$ data taken by the UA4 Collaboration [1]. A comparison of the $p\bar{p}$ and pp differential cross sections in the region of the dip can be found in Ref. [17].

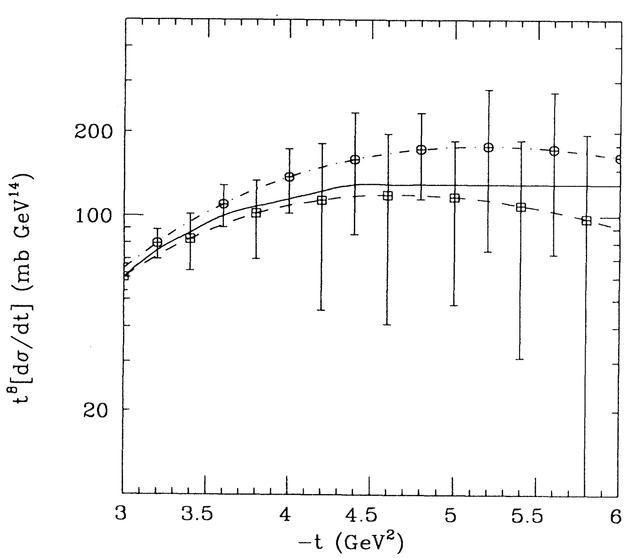

FIG. 1. A Plot of $t^8\frac{d\sigma}{dt}$ for pp elastic scattering at the CERN ISR. The squares and error bars correspond to the data at $\sqrt{s} = 23.4$ GeV, while the circles correspond to the data at $\sqrt{s} = 30.5$ GeV. The solid line is an average approximation to the data showing the approximate t^8 behavior of the differential cross section, characteristic of the Landshoff mechanism.

It is important to note that the $|t|$ range of these high energy collider data overlaps with the $|t|$ range of the low energy, large angle data, where striking structure has been observed in the two-spin observable A_{NN} [18] and where oscillations around an approximate s^{-10} falloff have been found in the 90° c.m. differential cross section [1]. Given the known s dependence of the three-gluon amplitude, we can definitely rule out the possibility that the Landshoff mechanism is involved in these low energy structures. Since this possibility has been widely considered, we will discuss the normalization issue here in some detail.

Since the Landshoff normalization factors $|L(t)|$ are dependent only upon $-t$, we can use existing cross section data at various s for comparable $-t$ values to phenomenologically determine their behavior. We can then extrapolate the Landshoff amplitudes back to the kinematic region where these t values correspond to known 90° c.m. data to find their effect on the oscillations of the cross section and A_{NN} spin asymmetry. For the purposes of this exercise, we have fit elastic pp data from the Argonne Zero Gradient Synchrotron (ZGS) and CERN (ISR) for various $\sqrt{s}$ ranging from 3 GeV/c to 62 GeV/c and for $-t$ from about 2 to 6 (GeV/c)^2. In analogy to previous analyses of the ISR data [1], we have constructed fits to the differential cross sections having the form

$$\frac{d\sigma}{dt}(\mu b/\mathrm{GeV}^2) = \beta \exp[-\delta|t|], \tag{2.10}$$

where β is an overall constant. When the differential cross sections are written in this form, we can extract the behavior of the Landshoff normalization amplitudes in a straightforward way.

The ANL data [19] at 90° c.m. covers a range of $\sqrt{s}$ from 3.3 to 5.1 GeV/c. There is a marked change in structure of the cross section for $-t$ near 7 (GeV/c)^2. Our fit to the data is

$$\begin{aligned}\frac{d\sigma}{dt}(\mu\mathrm{b}/\mathrm{GeV}^2) &= 6836\exp[-1.59|t|],\\ &\quad 3.8 \le -t \le 6.8(\mathrm{GeV}/c)^2\\ &= 24\exp[-0.76|t|],\\ &\quad 7.3 \le -t \le 11.3(\mathrm{GeV}/c)^2.\end{aligned} \tag{2.11}$$

The top line corresponds to the $-t$ region of our analysis.

Our fit to CERN PS data [20] at $p_0 = 19.2$ GeV/c, $\sqrt{s} = 6.15$ GeV/c, and large angle (60–90 degrees) gives

$$\begin{aligned}\frac{d\sigma}{dt}(\mu\mathrm{b}/\mathrm{GeV}^2) &= 38.5\exp[-1.65|t|],\\ &\quad 1.4 \le -t \le 5.0(\mathrm{GeV}/c)^2.\end{aligned} \tag{2.12}$$

The original analysis of the ISR data [1] was parametrized in a slightly different form and gave the differential cross section in (mb/GeV^2). Our analysis is equivalent, and we have written the cross section in (μb/GeV^2) for consistent comparison with other data over a similar $-t$ range. Our results at each $\sqrt{s}$ for 2.0 $\le -t \le 4.5$ GeV^2 are

$\sqrt{s}$(GeV)	δ(GeV^{-2})	β(μb/GeV^2)
23.4	1.71 ± 0.05	1.55 ± 0.20
30.5	1.54 ± 0.09	1.00 ± 0.25
44.6	1.71 ± 0.09	1.60 ± 0.40
52.8	1.77 ± 0.04	2.05 ± 0.20
62.1	1.77 ± 0.08	1.85 ± 0.40

(2.13)

which agree favorably with those in Ref. [1].

For comparison, we note that Donnachie and Landshoff [17] give the normalization of the cross section for the three-gluon exchange mechanism in the form

$$\left(\frac{d\sigma}{dt}\right) = (7\times 10^8)\alpha_s^6 t^{-8}\frac{P^4}{R^{12}}, \tag{2.14}$$

where P is related to the q-q scattering probability and R is an effective three-quark system radius. Our $|L(t)|^2$ factor therefore corresponds roughly to Landshoff's factor: $\alpha_s^6(P^4/R^{12})\times 10^8$. Note that the contribution of the three-gluon mechanism is dependent only on $-t$ and is independent of s, as is the experimental cross section at large s. This observation confirms our claim that, although the treatment of Ref. [17] applies to the region of fixed $-t$ and smaller angles, the strict $-t$ dependence of the normalization allows us to extrapolate the three-gluon contributions down to energies of the ANL at 90° c.m. for similar $-t$ values. We can then determine the relative importance of this mechanism to the elastic processes there.

We assume that the ISR data are in a region dominated by the Landshoff mechanism. Using Eqs. (2.11) and (2.13), we can calculate a ratio of the cross sections between the ISR and ANL data:

$$\left(\frac{d\sigma}{dt}\right)_{\mathrm{ISR}} \Big/ \left(\frac{d\sigma}{dt}\right)_{\mathrm{ANL}} = (2.4\times 10^{-4})\exp[-0.12|t|]. \tag{2.15}$$

When considering the effects of the Landshoff mechanism at 90° c.m., the u-channel terms become important, but are correctly accounted for in our full parametrization. Thus at ANL energies, the total Landshoff contribution is at most about four times the value in Eq. (2.15), or about 6×10^{-4} for this range of t values. Thus, although the Landshoff contribution to the ISR cross section may be large, its effect on the ANL 90° data is down by a factor of 10^{-4}. We conclude that another mechanism must be responsible for the structure of the data in this region. We now turn to an analysis of the amplitudes at 90° c.m.

III. STRUCTURE OF THE AMPLITUDES AT 90° c.m.

The normalization of the Landshoff mechanism given as determined from ISR data in the analysis of Sec. II rules out the possibility that this mechanism is instrumental in determining the "low energy" structure observed in the large angle cross section. There are two distinct types of phenomena which have been observed in the 90° c.m. cross section for pp elastic scattering and have presented theoretical challenges to our understanding: (1) oscillations in $s^{10}d\sigma/dt$ as a function of s and (2) structure in A_{NN} as a function of energy.

There have been attempts [9–12] to describe each of these effects in terms of an interference effect between a Landshoff amplitude and a quark-interchange amplitude. Based on the discussion in Sec. II, it appears that we must search elsewhere for the explanation of these unusual structures.

There are ample reasons to believe that a description of the amplitudes in terms of a constituent-based hard scattering mechanism in this kinematic region is sensible. As mentioned earlier, the cross section shows an approximate s^{-10} behavior, characteristic of a short distance process. As discussed by Brodsky, Carlson, and Lipkin and by Farrar, Gottlieb, Sivers, and Thomas [6], the most natural approach to these amplitudes is the quark-interchange model, or QIM. Using the symmetries of the minimal three-quark state in the proton and the conservation of quark helicity in the underlying hard scattering, the ratios of different pp-helicity amplitudes can be specified.

We would like to consider the possibility that the basic symmetries of the QIM model are reflected in the large angle amplitudes and that the structures that have proved so puzzling can be understood as corrections to these amplitudes. There are certain simplifications of the problem which occur because of the symmetries at 90° c.m. For example, since $t = u$ here, we have the kinematic constraints that

$$\Phi_4 = -\Phi_3, \quad \Phi_5 = 0. \tag{3.1}$$

In addition, we would like to implement the assumption that the double helicity-flip amplitude, Φ_2, can be neglected at the energies where the data exist. This imposes the constraint of hadronic helicity conservation at 90° c.m., which the data and our previous analyses have shown to be a reasonable approximation. Helicity conservation therefore implies that the observables at 90° can be understood in terms of only two independent amplitudes. For the observables we consider here, the expressions are simplifications of those written by Hendry [21]:

$$\Sigma = \frac{1}{2}[|\Phi_1|^2 + 2|\Phi_3|^2], \quad A_{NN} = \frac{2|\Phi_3|^2}{[|\Phi_1|^2 + 2|\Phi_3|^2]}. \tag{3.2}$$

We now implement the constraint that the amplitudes asymptotically approach the QIM result: $\Phi_1^Q = 2\Phi_3^Q$ at 90° c.m., to write

$$\Phi_1 = 2\Phi^Q + \hat{\Phi}_1, \quad \Phi_3 = \Phi^Q + \hat{\Phi}_3, \tag{3.3}$$

where Φ^Q is a smooth power-law behaved amplitude which characterizes the asymptotic observables, while $\hat{\Phi}_1$ and $\hat{\Phi}_3$ are subasymptotic corrections. If we write

$$\Sigma_0 \equiv \frac{1}{2}6|\Phi^Q|^2, \tag{3.4}$$

then

$$A_{NN}(\Sigma/\Sigma_0) = \frac{1}{3}\frac{|\Phi_3|^2}{|\Phi^Q|^2} = \frac{1}{3}\frac{|\Phi^Q + \hat{\Phi}_3|^2}{|\Phi^Q|^2}, \quad (1 - A_{NN})\Sigma/\Sigma_0 = \frac{1}{6}\frac{|\Phi_1|^2}{|\Phi^Q|^2} = \frac{2}{3}\frac{|\Phi^Q + \frac{1}{2}\hat{\Phi}_1|^2}{|\Phi^Q|^2}. \tag{3.5}$$

At this point it is convenient to factor out the overall phase of Φ^Q and write

$$A_{NN}\Sigma/\Sigma_0 = \frac{1}{3}|1 + \hat{\alpha}_3|^2, \quad (1 - A_{NN})\Sigma/\Sigma_0 = \frac{2}{3}\left|1 + \frac{1}{2}\hat{\alpha}_1\right|^2, \tag{3.6}$$

where $\hat{\alpha}_3 \equiv \frac{\hat{\Phi}_3}{\Phi^Q}$ and $\hat{\alpha}_1 \equiv \frac{\hat{\Phi}_1}{\Phi^Q}$. Using a smoothed version of the data, we have separately extracted

$$R_1 \equiv \frac{|\Phi_1|^2}{|2\Phi^Q|^2} = |1 + \frac{1}{2}\hat{\alpha}_1|^2, \quad R_3 \equiv \frac{|\Phi_3|^2}{|\Phi^Q|^2} = |1 + \hat{\alpha}_3|^2 \tag{3.7}$$

from these expressions and have shown them in Fig. 2. At this point, we note that our simplified approach to "amplitude analysis" has relied only on the symmetries present in the 90° amplitudes and the assumption of hadronic helicity conservation in order to separate observables depending on Φ_1 and Φ_3. The QIM constraint that $\Phi_1^Q = 2\Phi_3^Q$ at 90° enters indirectly into the coefficients in (3.6) and (3.7). If the data are correct, they show interference effects in both amplitudes with the structure in Φ_1 occurring at a lower energy than that in Φ_3.

The structure of A_{NN} for the range $4 < |t| < 7$ GeV^2 is relatively flat. The value of A_{NN} however, differs significantly from the QIM prediction of $\frac{1}{3}$. In this region, we find $|\Phi_1|^2 > |2\Phi^Q|^2$, while $|\Phi_3|^2 < |\Phi^Q|^2$. Whatever subasymptotic mechanism is invoked to explain these data, its impact on $|\Phi_1|$ and $|\Phi_3|$ can be seen directly in Fig. 2. The differential cross section in this region exhibits a relatively steady $-t^{-10}$ behavior. If we make the ansatz that $\Sigma_0 = 3.5 \times 10^8/(-t)^{10}$, then Σ has the same t behavior. This is a measure of the relative strength of the mechanism which "interferes" with the QIM amplitudes in this region to cause A_{NN} to dip below the QIM prediction. Over a considerable t range, the "corrections" are comparable to the QIM amplitudes.

In the region $|t| > 7$ GeV^2, the slope of the cross section changes and A_{NN} simultaneously begins to rise sharply. The differential cross section exhibits more of a t^{-8} be-

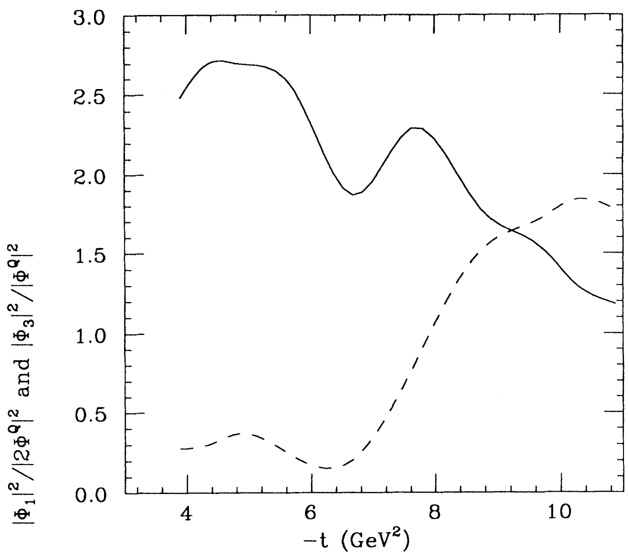

FIG. 2. Behavior of the ratios $R_1 \equiv \frac{|\Phi_1|^2}{|2\Phi^Q|^2}$ (solid curve) and $R_3 = \frac{|\Phi_3|^2}{|\Phi^Q|^2}$ (dashed curve) with $-t$, extracted from 90 degree c.m. data.

havior in this region. The amplitude Φ_3 now becomes more significant here, since $\frac{|\Phi_1|^2}{|\Phi_3|^2} \approx 1$ for $|t| \geq 8$ Gev^2. The mechanism responsible for the "oscillation" of the cross section and the rapid rise in A_{NN} now appears to enhance Φ_3 and suppress Φ_1 over their QIM values. It is possible to understand this switchover as an interference effect, where the phase and magnitude of the subasymptotic amplitudes are changing. Both $|\hat{\alpha}_1|$ and $|\hat{\alpha}_3|$ are of comparable magnitude here. Unfortunately, the data stop before we can confirm our hypothesis that $|\hat{\alpha}_1|$ and $|\hat{\alpha}_3|$ vanish at higher energies.

We have not discussed the specific dynamical mechanisms which may be responsible for these effects, except to rule out a significant contribution from the Landshoff three-gluon exchange diagrams discussed above. We note that Brodsky and deTeramond [22], have discussed the structure of A_{NN} in terms of a specific dibaryon resonance, which couples to Φ_3, but not to Φ_1. A dibaryon resonance with different quantum numbers may turn out to be the most economical explanation for the structure in Φ_1 at lower t. If we are to identify deviations of $\hat{\alpha}_1$ and $\hat{\alpha}_3$ from zero with nonpointlike scattering configurations, then the structure in Fig. 2 indicates that existing data are not in a regime where "nuclear transparency" should be a feature of $pA \to pp(A-1)$. This implies that it is important to push for experiments at higher energy on nuclear targets to see if nuclear transparency emerges in a regime where $\hat{\alpha}_1$ and $\hat{\alpha}_3 \to 0$.

The idea of Jain and Ralson [23] concerning the space-time structure of nuclear transparency imply that the subasymptotic corrections associated with $\hat{\alpha}_1$ and $\hat{\alpha}_3$ are suppressed in the nuclear environment. In their approach, the pointlike cross section, Σ_0 is a more useful quantity for understanding nuclear effects than is the physical cross section. Further discussion of this can be found in Ref. [23]. An important requirement for understanding these effects is to seek new data at higher energies. For many years, there have not existed experimental facilities with polarized proton beams to continue the experimental program responsible for these data. Fortunately, it is now possible to advocate new experiments to continue this study.

It is interesting that the quark interchange mechanism (QIM) is dominant over the three-gluon exchange diagrams in the region. This is confirmed by the fact that $\frac{d\sigma(p\bar{p})}{d\sigma(pp)} < \frac{1}{50}$ in this kinematic regime. This small ratio is consistent with quark interchange mechanisms. If Landshoff diagrams were important, the ratio would have to be much closer to unity. The evidence for contributions from nonleading twist or subasymptotic mechanisms which we have extracted from this analysis is quite indirect. Further experiments are necessary in order to confirm this basic approach. One place where subasymptotic dynamics is more accessible involves the measurement of the polarization asymmetry. We now turn to a brief discussion of the expectation for this observable.

IV. THE POLARIZATION ASYMMETRY AT LARGE ENERGIES

The expression of Eq. (2.3) for the polarization asymmetry can be written

$$P = \frac{-\mathrm{Im}[\Phi_5^*(\Phi_1+\Phi_2+\Phi_3+\Phi_4)]}{|\Phi_1|^2+|\Phi_2|^2+|\Phi_3|^2+|\Phi_4|^2+4|\Phi_5|^2}. \quad (4.1)$$

Our parametrization of the helicity-flip amplitude, Φ_5, in terms of the helicity-conserving amplitudes

$$\Phi_5 \equiv \frac{-\epsilon(-t)^{1/2}}{2m_p^2-t}[\Phi_1(s,t)+\Phi_3(s,t)+\Phi_4(s,t)] + \frac{-\epsilon(-u)^{1/2}}{2m_p^2-u}[-\Phi_1(s,t)+\Phi_3(s,t)+\Phi_4(s,t)] \quad (4.2)$$

is based on the idea of restructuring a proton wave function from scattered quarks which are approximately collinear. This approach builds in the constraint of hadronic helicity conservation so that at fixed angles, the polarization asymmetry vanishes: namely,

$$\lim_{\substack{s\to\infty \\ t/s \text{ fixed}}} P = \frac{c}{\sqrt{s}}. \quad (4.3)$$

Note that the form (4.2) for Φ_5 is also consistent with the kinematic constraints of (3.1) at 90° c.m.. However, at fixed t, the s dependence of Φ_5 in our approach becomes asymptotically equal to the s dependence of the helicity conserving amplitudes, since the form factor

$$f(t) \equiv \frac{-\epsilon(-t)^{1/2}}{2m_p^2-t} \quad (4.4)$$

is a function of t only.

As discussed in Refs. [9] and [13], this form factor represents the effect of a convolution over internal quark momenta with quark helicities conserved, but with the relative transverse momenta of the quarks within the proton taken into account. The $(-t)^{1/2}$ represents the $\sin(\theta/2)$ factor which must appear in a helicity flip amplitude, and the denominator is chosen to match dimensional counting rules. The complex parameter, ϵ, represents the overlap of a multiquark state with a proton. In an exact SU(6) model, ϵ would vanish. The approximate success of SU(6) for static properties of the proton suggests that ϵ should be small. The fact that Eq. (4.4) represents a convolution over phases makes it natural that ϵ should be complex. Our approach leads naturally to a polarization asymmetry which becomes independent of energy at fixed t.

It must be questioned whether this result is reasonable, given the known prediction from Regge theory that the polarization asymmetry should vanish asymptotically at fixed t. This Regge prediction is associated with the phase-energy relationship. For each Φ_i in (4.1) we can write

$$\mathrm{Re}[\Phi(s,t)] = \int_{\mathrm{RH}} \frac{\mathrm{Im}[\Phi_i(x,t)]}{x-s}dx + \int_{\mathrm{LH}} \frac{\mathrm{Im}[\Phi_i(x,t)]}{x-s}dx, \quad (4.5)$$

where the integrals are over the right-hand and left-hand cuts in the complex s plane. If asymptotically

$$\lim_{s\to+\infty} \mathrm{Im}[\Phi_i(s,t)] = C_i s^{\alpha_i},$$
$$\lim_{s\to-\infty} \mathrm{Im}[\Phi_i(s,t)] = \pm C_i s^{\alpha_i}, \quad (4.6)$$

Eq. (4.5) guarantees the Regge asymptotic phase satisfies

$$\lim_{s\to\infty} \frac{\mathrm{Re}[\Phi_i(s,t)]}{\mathrm{Im}[\Phi_i(s,t)]} \approx \frac{\cos(\pi\alpha_i)\mp 1}{\sin(\pi\alpha_i)}. \quad (4.7)$$

In the Regge region, factorizable Regge residues would require that Φ_5 would have the same asymptotic phase as Φ_1 and Φ_3, so that the bilinear in (4.1) would vanish. With a general multiquark exchange, however, there is no unique continuation from the left-hand cut to the right-hand cut, so there is no distinct phase factor for the different amplitudes. In this sense, it is quite reasonable for the parameter ϵ in (4.2) to be complex.

In our phenomenological approach, we can reproduce the existing polarization data [14] with a value

$$|\epsilon| \approx 0.15, \quad (4.8)$$

using (4.1) and (4.2). In this approximation, ϵ is predominately imaginary. The behavior of the polarization at large energies can therefore be summarized as follows:

(1) Small $|t|$. There is an overall $(-t)^{1/2}$ factor. In addition the polarization should fall off with energy, reflecting the coherent behavior of factorizable Regge poles.

(2) Large $|t|$. Outside the coherent region, the polarization becomes asymptotically energy independent. In order to be consistent with quark helicity conservation, the polarization should behave like $(-t)^{-1/2}$ as t is varied at large s.

Those polarization measurements which exist are consistent with these predicted regularities [24]. Further measurements of the polarization asymmetry at high energies are needed in order to test these underlying principles and to normalize the subasymptotic effects. There may be some unfortunate preconceptions about the importance of polarization asymmetry measurements which prevent their systematic study. While our basic approach of including helicity nonconserving effects with an empirical form factor is only an approximation, it provides an interesting guide to the type of information these measurements provide. The important thing is that helicity-nonconserving amplitudes are not zero, in contrast with the Regge prediction, but are small with a t dependence predicted from underlying principles.

The behavior of the polarization P (or A_N) as a function of $-t$ can be measured in elastic polarized pp scattering at Fermilab or the lower proposed energies, reached at the BNL Relativistic Heavy Ion Collider (RHIC), where the cross sections and luminosities would allow for a reasonable distinction from nonzero values. Specific predictions for values of P as a function of phenomenological parameters are given in Ref. [13]. The main emphasis here is that, if our assumptions about the helicity-nonconserving amplitudes are correct, we should see nonzero polarizations at large enough $-t$ values to be outside the coherent region.

V. SUMMARY AND CONCLUSIONS

Elastic scattering of hadrons has provided a wealth of information. Pursuing this information has been a process frought with false leads and complicated puzzles. For proton-proton elastic scattering, some regularities associated with the hard scattering region are beginning to emerge. Several independent phenomenological arguments support the identification of the t^{-8} structure observed in the data at $|t| \geq 4$ GeV^2 and $t/s \to 0$ with the three-quark scattering mechanism first proposed by Landshoff. There is some concern that the data are not in a truly "perturbative" regime since there is no evidence of a falloff associated with the running of the QCD coupling within the factor $\alpha_s^6(t)$ which appears in the cross section. Nor is there any evidence for an additional falloff associated with the Sudhakov form factor. There is no doubt that further study targeted at these effects is warranted and is necessary before we can be completely confident that a constituent-based mechanism can be isolated in the elastic scattering data.

In this paper, we have used the normalization of the Landshoff-model-helicity amplitudes which are taken from the phenomenological studies mentioned above and continued the amplitudes to smaller s values. There is little uncertainty in this exercise since the s dependence of the Landshoff amplitudes are well specified by the model. The continuation shows that the Landshoff amplitudes are too small (by 10^{-2} at the amplitude level) to be involved in the oscillations observed in $d\sigma/dt$ at 90° c.m. or in the sharp structure observed in A_{NN}. The most natural explanation of these striking phenomena involves the interference of some asymptotic mechanism with a dominant QIM model amplitude set. Brodsky and deTeramond have long advocated this type of explanation for the structure of A_{NN}. They propose a dibaryon resonance in Φ_3 associated with the opening of the $pp \to ppD\bar{D}$ threshold. The definitive test of this proposal involves measurements in other inelastic channels.

Using the symmetries in the amplitudes at 90° c.m. along with the assumption that the double helicity-flip amplitude can be neglected, we have expressed the data in terms of Φ_1 and Φ_3. Using the QIM relation $\Phi_1 = 2\Phi_3$ as a starting point, we find the data suggest a structure which interferes with the dominant QIM amplitudes. This interference occurs at a lower energy in Φ_1 than in Φ_3, but is approximately the same magnitude in each amplitude. The data disappear in a kinematic regime where there is a lot of structure, and it would be interesting to continue these measurements at some higher energies. There exists a real opportunity to do these measurements at the Brookhaven Alternate Gradient Synchrotron (AGS) with the addition of a partial Siberian snake to allow for polarized beams. It will be interesting to see whether the data approach the value of $A_{NN} = \frac{1}{3}$ as predicted by the QIM model and whether the cross section oscillations fade away at higher energy.

Another type of measurement which can provide new insight involves the single-spin polarization asymmetry. Using factorizable Regge poles, the asymmetry should fall with energy at fixed $-t$. However, outside of the coherent Regge region, dynamical mechanisms based upon multiple quark scattering, such as the Landshoff or QIM mechanisms, allow for a nontrivial phase between different amplitudes. Hadronic helicity conservation implies that the amplitude Φ_5 should have a different power behavior than the helicity-conserving amplitudes. However, this can be accommodated with a rather mild $t^{-1/2}$ falloff of the polarization observable. Existing data are consistent with this behavior but do not provide a stringent test of the basic idea. Our simple model relates different amplitudes using a spin-flip "form factor" involving a small parameter related to SU(6) breaking.

Finally, within the context of our model, we have looked at the data for evidence of nontrivial behavior of the observables associated with the running of the QCD coupling or with the falloff of Sudhakov form factor. We could find no evidence of these effects. This may indicate that while constituent based models can provide important insight into the structure of $pp \to pp$ amplitudes, the data are not in a kinematic regime where perturbative calculations can be attempted.

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[1] A. Bohm *et al.*, Phys. Lett. **49B**, 491 (1974); E. Nagy *et al.*, Nucl. Phys. **B150**, 221 (1979); Foley *et al.*, Phys. Rev. Lett. **15**, 45 (1965).
[2] S. J. Brodsky and G. R. Farrar, Phys. Rev. Lett. **31**, 1153 (1973); Phys. Rev. D **11**, 1309 (1975).
[3] V. A. Matveev, R. M. Muradyan, and A. V. Tavkheldize, Nuovo Cimento **7**, 719 (1973).
[4] G. R. Farrar, H. Zhang, A. A. Globlin, and I. R. Zhitnitsky, Nucl. Phys. **B311**, 585 (1989); G. R. Farrar, E. Maina, and F. Neri, Phys. Rev. Lett. **53**, 28 (1984); **53**, 742 (1984).
[5] J. F. Gunion, D. Millers, and K. Sparks, Phys. Rev. D **33**, 689 (1986); D. Millers and J. F. Gunion, *ibid.* **34**, 2657 (1986).
[6] G. R. Farrar *et al.*, Phys. Rev. D **20**, 202 (1979); S. J. Brodsky *et al.*, *ibid.* **20**, 2278 (1979).
[7] M. Jacob and G. C. Wick, Ann. Phys. (N.Y.) **7**, 404 (1959).
[8] P. V. Landshoff, Phys. Rev. D **10**, 1024 (1974); P. V. Landshoff and D. J. Pritchard, Z. Phys. C **6**, 69 (1980).
[9] G. P. Ramsey and D. Sivers, Phys. Rev. D **45**, 79 (1992).
[10] C. E. Carlson, M. Chachkhunashvili, and F. Myher, Phys. Rev. D **46**, 2891 (1992).
[11] B. Pire and J. P. Ralston, Phys. Lett. **117B**, 233 (1982).
[12] J. Botts, Nucl. Phys. **B353**, 20 (1991).
[13] G. P. Ramsey and D. Sivers, Phys. Rev. D **47**, 93 (1993).
[14] G. W. Abshire *et al.*, Phys. Rev. Lett. **32**, 1261 (1974).
[15] V. Sudhakov, Zh. Eksp. Theor. Phis. **30**, 87 (1956) [Sov. Phys. JETP **3**, 65 (1965)]; J. Botts and G. Sterman, Phys. Lett. B **224**, 201 (1989); Nucl. Phys. **B325**, 62 (1989).
[16] G. P. Lepage and S. J. Brodsky, Phys. Rev. D **22**, 2157 (1980).
[17] A. Donnachie and P. V. Landshoff., Z. Phys. C **2**, 55 (1979); Phys. Lett. **123B**, 345 (1983); Nucl. Phys. **B231**, 189 (1983); **B244**, 322 (1984); **B267**, 690 (1986).
[18] A. Lin *et al.*, Phys. Lett. **74B**, 273 (1978); E. A. Crosbie *et al.*, Phys. Rev. D **23**, 600 (1981).
[19] C. W. Akerlof *et al.*, Phys. Rev. **159**, 1138 (1967).
[20] J. V. Allaby *et al.*, Phys. Lett. **28B**, 69 (1968).
[21] A. W. Hendry, Phys. Rev. D **10**, 2300 (1974).
[22] S. J. Brodsky and G. deTeramond, Phys. Rev. Lett. **60**, 1924 (1988).
[23] P. Jain and J. P. Ralston, in *Proceedings of the International Conference on Elastic and Diffractive Scattering*, Providence, Rhode Island, 1993 (unpublished).
[24] See, for example, M. Block, K. Kang, and A. White, Int. J. Mod. Phys. A **7**, 4449 (1992).