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Diffraction-Dissociation Model of High-Energy Nucleon-Nucleon Interactions*

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A model, discussed previously, where high-energy nucleon-nucleon interactions are considered from the view that diffraction-dissociation processes may be dominant at high energies, is used to calculate production multiplicities and meson and nucleon momentum spectra for proton-proton interactions at laboratory energies between 30 and 1500 GeV.

INTRODUCTION

In a previous paper,[1] referred to here as RA1, we considered the possibility that diffraction-dissociation processes[2] may dominate the very high-energy interactions of hadrons. We developed a specific, simple model of such processes and showed that the results concluded from the model were consistent with the information then available concerning very high-energy nucleon-nucleon interactions. This information was derived almost wholly from studies of cosmic-ray interactions and the diffusion of cosmic rays through the atmosphere.

Now that more detailed and more precise measurements of high-energy interactions will be available from measurements at the accelerator now nearing completion at the National Accelerator Laboratory and with the intersecting storage rings constructed at CERN, it seems desirable to use the model to compute properties of proton-proton interactions which might be easily measured using these facilities.

We first review the model with the particular purpose of defining clearly the assumptions made and the specific consequences of these assumptions.

FIRST ASSUMPTION

We assume that at sufficiently high energy each interaction of two nucleons results in the production of two separate, noninteracting states. Each state is identified with one of the initial nucleons inasmuch as the four-momentum difference between the initial nucleon and the final state associated with that nucleon is small. Further, the final state has the same quantum numbers, excepting spin and parity, as the initial nucleon. Each of the final states resulting from a proton-proton interaction will have a charge of 1, a hypercharge of 1, isotopic spin $\frac{1}{2}$, T_3 equal to $\frac{1}{2}$, and the state will have the same properties with respect to SU_3 as the proton member of the baryon octet.

We find it convenient to refer to these final states as "fireballs" inasmuch as they are closely related conceptually to hypothetical states so named which have been introduced[3] to account for phenomena observed in the high-energy interactions of cosmic-ray hadrons. However, we note a difference between our fireballs and the traditional fireballs. Our fireballs are so defined that their baryon number is the same as the associated incident particle: If the associated particle is a proton, the baryon number of the fireball will be 1. In the traditional view of the same reaction, the whole final state is divided into two parts where one part is a final nucleon or nucleon isobar state and the other part, with baryon number 0, is defined as the fireball. If the velocities of the two parts are nearly the same, the definitions tend to be equivalent.

This assumption of the diffraction-dissociation production of two final-state fireballs represents an extrapolation of the knowledge we have of exchange interactions at moderate energies (below 30 GeV). We observe that the total cross sections for those processes in which the quantum numbers of the states do not change – or the exchanged particle, the Pomeranchukon, has the quantum numbers of the vacuum – are nearly independent of energy while all other exchange interactions fall off rapidly with energy. If the exchanged quantum numbers correspond to the quantum numbers of a nonexotic particle (a meson which can be constructed from a quark and antiquark or a baryon which can be constructed from three quarks) the cross section appears to fall off with energy as s^{-x}, where s is the square of the energy in the center-of-mass system and x is greater than 0. Those interactions which cannot be described in terms of the exchange of a nonexotic particle or the vacuum appear to fall off with energy even more quickly. It then seems plausible that it might be possible to classify all interactions in terms of the properties which are exchanged, and only those reactions where the vacuum is exchanged will be important at very high energies.

This assumption, alone, leads to some interesting consequences, and some of these results might be easily observed and serve as a test of the assumption.

(1) Since each state produced in a proton-proton interaction will have a charge of 1, each of the states must decay to an odd number of charged particles.

(2) Since each state must have the same isotopic spin as the proton, charge independence leads to relations in the intensity of decay products. The most interesting and easily accessible prediction concerns the intensities of different meson charge states. At any energy and at any angle the intensity of neutral pions will be equal to one-half of the intensity of charged pions.

(3) Each final state must transform with respect to SU_3 as the proton member of the baryon octet. This leads to various relations which would be valid in the limit of complete SU_3 symmetry. Since the symmetry is rather badly broken, the relation of symmetry predictions to nature is always somewhat obscure. However, we present one example of such a prediction noting that the required measurements seem to be quite difficult. There is a relation in the intensity of Δ^-, Y^{*-}, Ξ^{*-}, and Ω^- such that the intensity can be expressed in the form

$$I = a + bY ,$$

where a and b are constants and Y is the hypercharge of the member of the quartet. In the limit of SU_3 symmetry, this relation would be valid at any energy and angle.

SECOND ASSUMPTION

In order to consider further characteristics of high-energy interactions, it is necessary to introduce dynamic assumptions. We make some qualitative comments before listing the conclusions derived in RA1. It is plausible that the only exchange reactions which are important are those which proceed with small values of four-momentum transfer. Since the four-momentum transfer required to produce a pair of states at a given center-of-mass energy increases with the mass of the states, and the four-momentum to produce states of a definite invariant mass decreases with increasing energy, we can consider that a definite threshold exists for the production of a pair of states of a given mass. Then as the energy increases, new states will become energetically available for production through the diffraction-dissociation mechanism. Since the total cross section seems to be constant, the additional cross section for the production of more massive states must be compensated by a decrease in the cross section for production of the lighter states. If the interaction is factorizable, that is, if the probability of the production of a state of mass M_f through the dissociation of the forward proton is independent of the probability of the production of a state of mass M_b through the dissociation of the backward proton, this compensation leads to a definite relation between the mass spectrum of the states produced through the dissociation process and the rate of decrease of the cross section for the production of states of specific masses with energy.

In particular, we assume that the cross section for the production of a pair of states j is not, then, constant but varies inversely with the logarithm of the energy as

$$\sigma_j = \frac{g_j^2}{2\alpha' \ln(s/s_j)} , \tag{1}$$

where s is the square of the center-of-mass energy and $g_j^2/2\alpha'$ is a constant written in this form for convenient comparison with other formulas. As with other two-body processes, we can presume that the cross section varies with t, the square of the four-momentum transfer, as

$$d\sigma_j/dt = A \exp Bt$$

and

$$\sigma_j = \int_{-\infty}^{t''} (d\sigma_j/dt)\, dt = (A/B) \exp Bt'' , \tag{2}$$

where A and B are appropriate constants and t'' is the maximum value of t allowed kinematically for the reaction j at the squared energy s. Clearly, the relations (1) and (2) are in accord if we set $A = g_j^2$ and $B = 2\alpha' \ln(s/s_j)$, where the parameters g_j^2, α', and s_j may, or may not, be associated with similar parameters which are used in the Regge description of particle reactions. The partial cross section for the reaction j then decreases with energy inversely with the energy for all s sufficiently large that t'' is nearly zero. This result, well known in its application to elastic scattering, applies to all cross sections in this diffraction-dissociation model of reactions.

Equation (1) is only consistent with the constraint that the total cross section for the production of all states is independent of energy if the decrease in cross section with respect to squared energy s of those states such that $-t''$ is near zero is precisely compensated by the increase in cross section with energy for those states with invariant masses so large that $-t''$ is large, noting, again, that for any final states of definite invariant mass squared $-t''$ decreases with increasing energy squared s. This precise balance requires a definite spectrum of masses for the states produced in the interaction. We write in general

$$d\sigma(M_a, M_b) = S(s)\, dM_a dM_b F(M_a) F(M_b) \times \exp[2\alpha' \ln(s/s_{ab})]t' , \tag{3}$$

where

$$S(s) = \sigma_j/[2\alpha' \ln(s/s_j)] \approx \sigma_{tot}/[2\alpha' \ln(s/s_0)] ,$$

where σ_{tot} is the total cross section and s_0 is an appropriate average of the scale constants s_j. The invariant masses of the forward and backward states are M_a and M_b, and t' is the square of the four-momentum transfer for production of the states at 0°. This expression is consistent with the relation of Eq. (1) for the differential cross section which has the characteristic diffraction-scattering shape. We note that the interaction is factorizable in the approximation that the values of the scale factors s_j are independent of the masses M_a or M_b (or if s_{ab} has a special form such as $s_{ab} = M_a M_b$) inasmuch as the probability of producing a state of a definite mass M_a' is independent of the accompanying mass M_b except through the constraints on the four-momentum transfer t'.

If such a factorization condition holds, we were able to show in RA1, that for large masses, at very large energies where the constraints on the four-momentum transfer can be ignored, $F(M)$ has the asymptotic form

$$F(M) \rightarrow [M(\ln M)^{1/2}]^{-1} \tag{4}$$

independent of the values of α' and s_0.

At moderate energies which are accessible to experiment, the character of the spectrum will depend upon the values of the parameters α' and s_0. The measurements of $d\sigma_j/dt$ for various states by Anderson *et al.*[4] at a laboratory energy near 30 GeV suggest that the differential cross sections for diffraction dissociation to many states might be adequately described by a form like Eq. (1) if α' is taken as equal to $1/M_n^2$, the characteristic slope for Regge trajectories, and s_j is then about 7 GeV^2. Using this value for all s_j and the constraint that the cross section is independent of energy, the form of $F(M)$ is defined by Eq. (3). While we were unable to find an analytic solution for $F(M)$, we were able to determine values of $F(M)$ numerically, using a computer, such that Eq. (3) was satisfied and the total cross section was constant from 15 to 50 000 GeV. For the purpose of exposition, we determined an analytic approximation to the numerical results:

$$F(M) = KA[A^2 + (M/M_n - 2)^2 \ln(M/M_n)]^{-1/2} ,$$

where (5)

$$A = M_n/M + 0.6 \quad \text{and} \quad M > M_n .$$

The value of the constant K was set so the integral of the differential cross section of Eq. (3) was equal to the measured total cross section.

Since the lowest possible mass for a fireball is the nucleon mass, and the next lowest mass must be a mass as large as the mass of a nucleon plus a pion, some specific modifications must be made to the spectral function of Eq. (5) at such low masses. Noting that the diffraction production of states with a mass less than 1.35 GeV/c^2 seems to be unimportant,[4] we set the contribution to states with a smaller mass to zero and, rather arbitrarily, give the singular nucleon state a weight equal to one-half of the amount subtracted.

The resulting fireball mass spectrum is shown in Fig. 1 for fireballs produced by proton-proton interactions of various energies. The asymptotic spectrum, which is shown with an arbitrary nor-

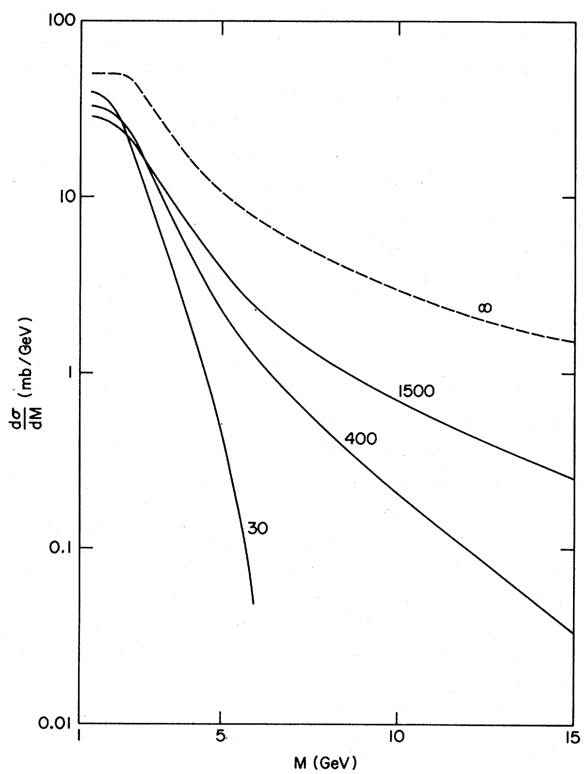

FIG. 1. The theoretical differential cross section for the production of fireballs in proton-proton interactions is plotted as a function of fireball mass, M, for various energies of the incident proton in the laboratory system. The asymptotic spectrum is presented on an arbitrary scale.

malization, has the form of Eq. (5). The difference between the asymptotic spectra and the spectra at finite energies follows from the restrictions on the production of high-mass states which follow from the conservation laws and the limitations on large four-momentum transfers.

We note here that the elastic scattering so calculated, where the final state of each incident nucleon has the mass of a nucleon, will be small (e.g., 0.5 mb at 300 GeV) compared to the observed elastic scattering cross section. We suggest that the cross section which we calculate corresponds to the square of the real part of the elastic scattering amplitude while the imaginary part, which is much larger, results from the requirements of the absorptive parts we are discussing together with the unitarity of the scattering amplitude. We consider that the diffractive-dissociation processes at high energy are, in some sense, weak processes and then the amplitudes for these reactions are nearly real, and then it is the real part of the elastic scattering amplitude which should be closely related to the reaction amplitudes. The increasing sharpness with energy of the angular distributions for the diffraction-dissociation reactions, as expressed by Eq. (1), is consistent with this view as this behavior indicates a physical situation where the interaction area increases with energy becoming much larger than the total cross section, and the eikonal interaction probability becomes small.

The second assumption can then be summarized as: *The mass spectra of the final states is expressed by the application of Eqs. (2) and (4).*

There are some specific consequences of this assumption though not all are easily accessible to experiment.

(4) The invariant-mass spectrum can be measured in principle and should follow the predictions defined in the assumption except, possibly, for small masses where the specific character of individual states may be much more important than the averages treated here.

(5) Certain factorization properties should obtain. The low-mass states of the backwards fireball might be easily observed, and the character of the production and decay of these states should be almost independent of the energy of the incident proton.

THIRD ASSUMPTION

In this model of high-energy interactions, many properties of the interactions will be determined by the character of the decays of the final-state fireballs. To consider the decays of the fireballs we must consider characteristics of the fireballs which are not defined by our first assumption. In particular, we will assume that the fireball can be described (approximately) as a quasistationary state and consider the consequences of this assumption. By a quasistationary state, we mean a state which has a lifetime longer than characteristic interaction times and a definite angular momentum and parity. The particles emitted by such a state will be emitted with a distribution in the rest mass system of the fireball which is fore-and-aft symmetric with respect to the direction of motion of the fireball. Indeed the decay distributions for any kind of particle will not likely be very different from isotropy and will not, in any case, be more singular than a distribution corresponding to classical isotropy in two dimensions where the particles are emitted in a plane perpendicular to the angular momentum axis. The distribution might then be expressed as

$$dN/d\Omega = A + B/\sin\theta ,$$

where θ is the direction of the emitted particle with respect to the direction of motion of the fireball. If the spin of the fireball is small compared to the number of particles emitted, we might expect that $A \gg B$. If the spin is large compared to the number of emitted particles, $B \gg A$.

In fact, we must expect that a final fireball state of definite invariant mass must represent a superposition of states of different parities and angular momenta. We then assume, implicitly, that over a sufficient interval of invariant mass, the contributions of the various states of different angular momentum and parity will be largely incoherent and that the decay distributions will not exhibit any large, systematic, variations from the symmetries expected from single states.

If we accept these assumptions, we can determine the momentum distribution of the fireball decay products in the center-of-mass system of the fireball by considering the experimentally determined transverse-momentum distribution of particles from high-energy interactions. It is well known[5] that the intensity of pions produced in high-energy hadron interactions varies with transverse momentum, p_t, approximately as

$$dN/d\Omega\, dp_t \propto \exp(-p_t/a) ,$$

where the mean transverse momentum is about 0.400 GeV/c. If we make the approximation that the pions are emitted from the fireball isotropically, we find that a momentum distribution of the form

$$dN/dp \propto p^2 \exp(-p/a') ,$$

where (6)

$$a' = 0.15 \text{ GeV}/c$$

results in a laboratory distribution of transverse momentum in approximate accord with observations. The mean transverse momentum of the pions in the laboratory system will be larger than in the fireball center-of-mass system as a result of the added effects of the fireball transverse momentum from production distributions and recoil effects, so that in the laboratory the effective value will be near 0.20 GeV/c. The transverse-momentum spectra from this model can then be expressed approximately as

$$dN/dp_t \propto p_t \exp(-p_t/a),$$

where (7)

$$a = 0.20 \text{ GeV}/c.$$

While the particular form for the momentum distribution expressed in Eq. (7) seems indicated by the experimental measurements of the angular distributions of mesons produced at high energies, the results which we will present do not depend upon the detailed character of this form, but are sensitive only to the mean emission energy of the pion which will be about 0.4 GeV. Indeed, a computationally simpler form was used in RA1 which was adequate for the averages considered in that analysis.

While the distribution of the transverse momentum of other particles produced in high-energy hadron-hadron interactions is not much different than the distributions for pions, we do not consider the production of other particles at this time making the implicit assumption that pion production dominates at the interaction energies we consider and that the production of other particles does not substantially affect the specific conclusions we reach.

Although the choice of a momentum distribution and angular distribution for the pions emitted from a fireball largely defines the character of fireball decay, the final distributions will be affected somewhat by more specific details of the decay. In particular, the fluctuations in decay properties which are so important in the analysis of the decays of single events are affected by detailed decisions concerning the character of the decays. For this reason, and for the sake of completeness, we explain in some detail the calculational procedures used to determine the results which are presented in the next section.

The results were calculated using a computer (the Brookhaven National Laboratory CDC 6600 computer) and Monte Carlo-type procedures. For each of a large number of events, the masses of the two fireballs were determined randomly following the prescriptions of Eqs. (3) and (5). The fireballs were assumed implicitly to decay by emitting pions consecutively. Each pion was presumed to be emitted from the fireball, in the fireball center-of-mass system, in a random direction with a random momentum fitted to a probability distribution of the form of Eq. (6). The charge of the pion was set randomly as positive, negative, or neutral with equal probability. After each such emission of a pion, the pion momentum and the fireball recoil momentum were calculated, and the fireball was then allowed to emit another pion with random momentum, direction and charge, again, in the center-of-mass system of the fireball. This procedure was modified when the decay would violate conservation laws by leaving a baryon state with a mass inconsistent with the remaining charge. Then the decay was "canceled," and the fireball was allowed to break up into two particles or three particles as necessary to conserve charge and energy. The momentum distribution for three-particle decays was determined by assuming that the momentum correlations of the final-state particles would follow the distributions defined by relativistic three-body phase space (equal areas on a Dalitz plot). When the decay of the remanent fireball state by a pion of a specific charge was forbidden, inasmuch as the remaining state would have a mass and charge incommensurate with decays to nucleons and pions, the sign of the charge was reversed. In this way the ratio of 2 : 1 for the decay of charged to neutral pions, which is required for the decay of isotopic-spin-$\frac{1}{2}$ states, was retained.

In this description of fireball decays, which considers the fireball as a baryon which decays to the nucleon ground state by consecutive emission of mesons even as an excited atom returns to the ground state by consecutive emission of photons, the nucleon retains a final recoil momentum which is of the same magnitude, but a little larger, than the mean momentum of the mesons. Therefore, the distribution of transverse momentum of the nucleons will be similar to that of the mesons.

We noted previously (RA1) that the results of this description of fireball decays, where the pions are emitted consecutively with a random momentum distribution of the type defined by Eq. (6), result in momentum distributions which are very similar to that derived from the assumption that decay probabilities were determined solely by volume in relativistically invariant phase space. We note that such phase-space distributions result in an absence of directional correlations between the directions of decay of the particles. So it is not wholly surprising that a different description of such decays, which minimizes the directional correlations between the different particles, may result in momentum distributions which are nearly the same as those which result from the assumption that the decay distributions are dominated by the available final-state phase-space volumes. While this conclusion might be of some intrinsic interest, suggesting that complex reactions can result in phase-space-like distributions though no thermal equilibrium is involved, the result is important to this description of the decay of fireballs inasmuch as it suggests that no important results would be much changed if a thermodynamic description of fireball decay were used rather than the "consecutive meson emission" model used here.

In summary, the third assumption, concerning the character of the decay of the fireballs, can be stated as: *We assume that the fireball may be described (approximately) as a quasistationary state which decays emitting mesons with a random momentum such that the distribution of magnitude of the momenta is described by Eq. (6).*

Some important consequences of this assumption must be considered rather as input data rather than predictions, since the properties in question have long been known and have led implicitly, to the selection of the model and the construction of the assumption. In particular:

(6) The transverse-momentum distribution of particles produced in high-energy hadron-hadron interactions will be nearly independent of the longitudinal momentum of the particle and of the center-of-mass energy of the interaction.

Less obviously, if the model is valid:

(7) The decay products which can be assigned to a forwards or to a backwards direction must have momenta which are consistent with their emission from a common center with a mean momentum of about 0.40 GeV/c. The angular distribution of the particles so emitted should not differ greatly from isotropy.

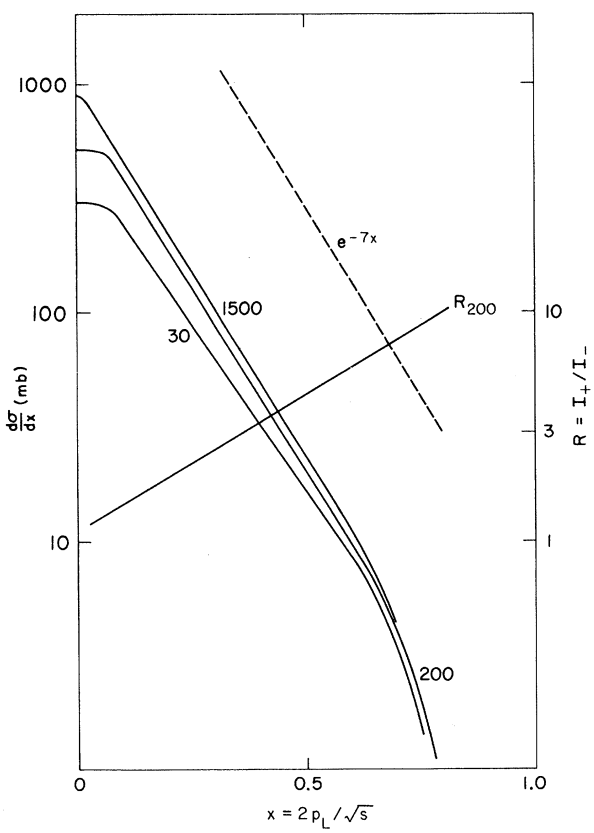

FIG. 2. The predicted differential cross section for meson production as a function of x, the ratio of the longitudinal meson momentum in the center-of-mass system to the maximum possible momentum, is plotted for various incident energies in the laboratory system. The ratio of positive to negative mesons produced by the interaction of protons with 200 GeV in the laboratory system is also shown as a function of x.

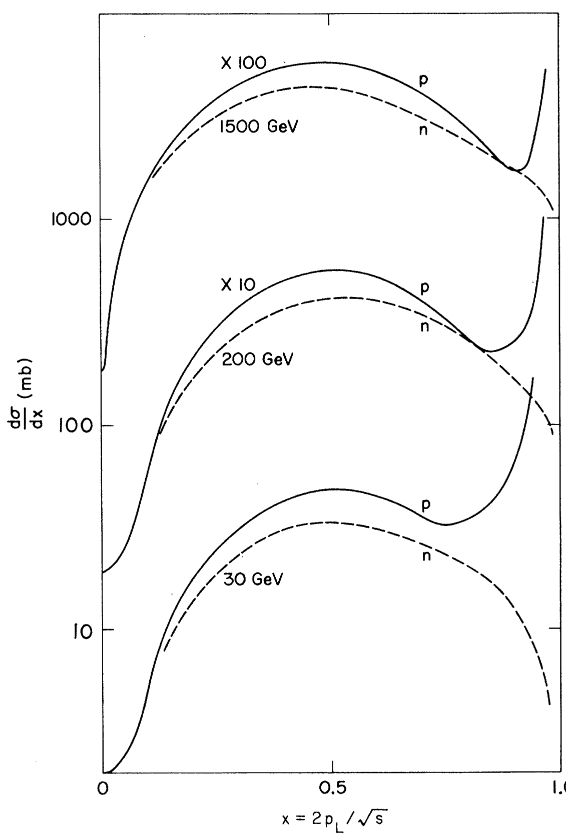

FIG. 3. The theoretical differential cross section for nucleons emitted in inelastic proton-proton interactions is plotted as a function of x, the ratio of the longitudinal momentum of the nucleon in the center-of-mass system to the beam momentum. The results are shown for proton and neutron production in the forward direction for various beam energies measured in the laboratory system.

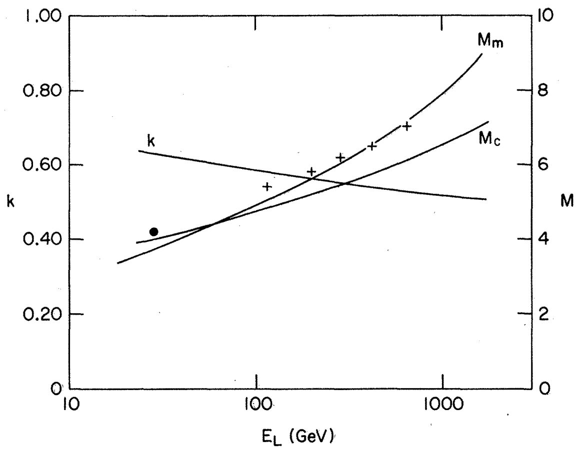

FIG. 4. The calculated multiplicity of mesons, M_m, and the multiplicity of charged particles, M_c, produced in the inelastic interactions of protons with protons are presented as a function of the energy of the incident proton in the laboratory system. Also shown is the predicted mean nucleon elasticity k as a function of incident proton energy where the elasticity is the ratio of the energy of the most energetic nucleon traveling in a specific direction (e.g., forwards) in the center-of-mass system to the beam energy in that system. The laboratory inelasticities are essentially the same.

FURTHER PREDICTIONS OF THE MODEL

Excepting the production of strange particles and baryon-antibaryon pairs, the model of high-energy nucleon-nucleon interactions introduced here is an almost complete description of such interactions. We present here further predictions of the model of proton-proton interactions emphasizing characteristic results which might be easily accessible experimentally. The results are derived from Monte Carlo calculations according to the prescriptions presented in the previous sections. As is often the case with such numerical calculations, the calculations do not always provide much insight into the analytical or physical bases of the ensuing results. Many of the predictions are presented for interactions at different laboratory energies ranging from 30 GeV (the maximum energies at the Brookhaven AGS and at the CERN PS) to 1500 GeV (relevant to measurements at the intersecting storage rings at CERN). However, we do not believe that the model is truly relevant at energies so low as 30 GeV, and the predictions presented for this energy are not expected to be reliable. The predictions:

(8) Figure 2 shows the meson momentum spectra in the center-of-mass system where the intensity is presented as a function of the longitudinal momentum. It is convenient to express the longitudinal momentum in terms of the ratio, $x = p_L/p_{max}$, where p_L and p_{max} are the longitudinal momentum and the maximum momentum allowed by the conservation laws: $p_{max} = \frac{1}{2}s^{1/2}$. Only positive values of x are plotted on the figure.

These curves present the spectra for all pions. The ratio of positive to negative pions at various longitudinal momenta for pions produced by 200-GeV proton-proton interactions is also shown. Expressed in terms of x, the variation of the ratio is almost the same for other energies. At all values of x, the neutral mesons constitute one-third of the total meson flux.

(9) Figure 3 shows the nucleon intensity as a function of x in the center-of-mass system. Note that only positive values of x are plotted, and then the figure refers to the nucleon intensity in one direction.

(10) The nucleon elasticity in the laboratory system, k, is shown in Fig. 4. The inelasticity is defined as the ratio of the average energy of the leading (that is most energetic) nucleon to the beam energy. The calculated values are in accord with values deduced from the analysis of cosmic-ray interactions. The asymptotic value – at infinite energy – of the inelasticity is zero on this model, but this value is reached only very slowly as shown in RA1 (where the inelasticity, $1-k$, is plotted).

(11) The mean meson production multiplicity, M_m, and the mean number of charged particles emitted in an interaction, M_c, are plotted as a function of laboratory energy in Fig. 4. The averages are calculated excluding elastic scattering. The crosses show the charged-particle multiplicities observed in the interactions of the charged hadrons in the cosmic-ray flux with protons,[6] while the solid circle shows the charged-particle multiplicity determined at accelerator energies.[7] According to this model, the asymptotic multiplicity will vary approximately as $(\ln s)^{1/2}$ for either particle production or charged-particle emission.

(12) The model also predicts multiplicity distributions. The graph of Fig. 5 shows the distribution of multiplicities of charged-particle emission for various laboratory energies. The dashed line connects points showing experimental values at an average energy of 300 GeV. While the agreement between experiment and theory is not very good, we note the possibility of appreciable biases in this difficult cosmic-ray measurement. Also, the importance of discrete production states and particular final-state interactions, necessarily neglected in a treatment concerned with broad averages, can seriously distort the theoretical predictions at small multiplicities.

(13) The assumption of factorization demands a certain lack of correlation between the multiplici-

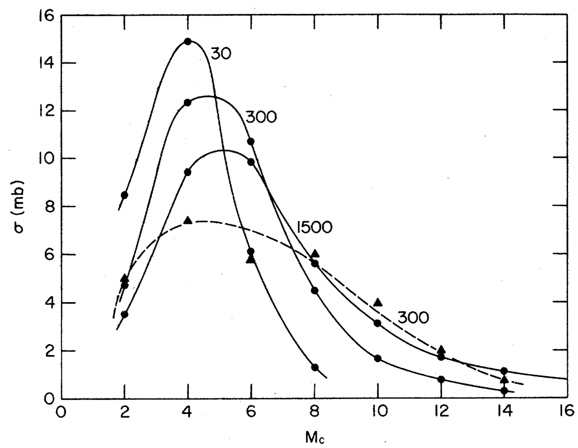

FIG. 5. The circles show predicted cross sections for the emission of various numbers of charged particles in inelastic proton-proton interactions. Theoretical results are plotted for various incident proton energies. The triangles show experimentally observed multiplicities from the interaction of cosmic-ray hadrons with protons. The lines are drawn as visual aids.

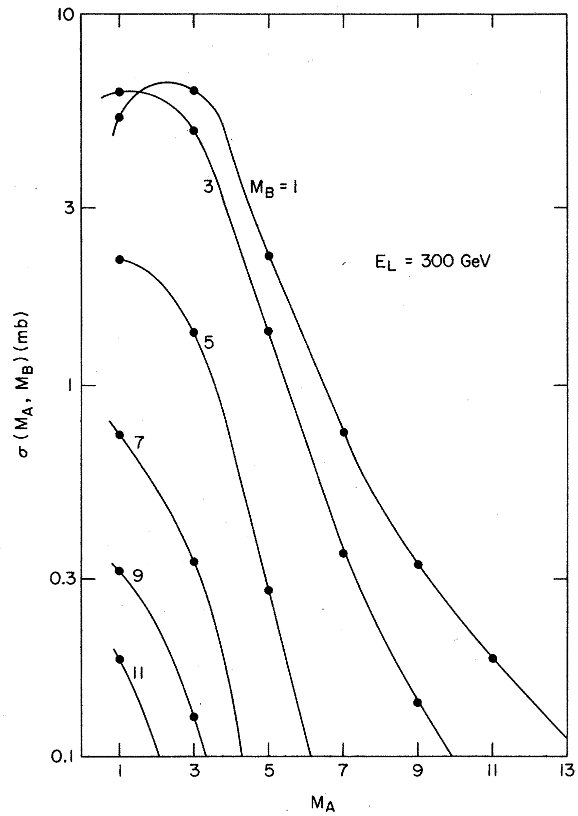

FIG. 6. The circles show the predicted cross sections for the production of M_A charged particles in the forward direction in the center-of-mass system and M_B particles in the backwards direction for proton-proton interactions where the laboratory energy of the incident proton is 300 GeV. The lines are drawn as visual aids.

ties derived from the disassociation of the forward fireball and the multiplicities from the breakup of the backward nucleon. Some of the effects of this factorization hypothesis are shown in the plots of Fig. 6, where the cross sections for different charged-particle emissions in the forward direction are plotted for different backward charge multiplicities. Again, the effects of discrete states, not considered in the theory which is concerned with broader averages, may modify the precise character of the distributions, but if factorization is a good approximation, evidence for its validity – or failure – should be exhibited in such plots.

COMMENTS AND CONCLUSIONS

We believe that this model of high-energy processes has, at least, the virtue of definiteness and accessibility – and then vulnerability – to experimental measurements. The only arbitrary numerical choice was the choice of magnitude of the contribution from the singular nucleon state to the fireball mass spectrum. Nothing would change very much if a smaller contribution of this state, for example, zero, were predicated but a much larger contribution leads to somewhat smaller multiplicities and a shift of the multiplicity spectrum, as shown in Fig. 5, to smaller values.

Of course, there are undoubtedly numerous alternatives to the logical choices made in the construction of the theory, even within the general theme. These choices cannot be enumerated very well, however, because for psychological reasons the choices cannot usually be remembered: The choices which are made always seem, eventually, unique and inevitable. Our choice of factorization is an exception to this rule, however. We rather doubt that factorization of the sort we use will be fundamental, but we believe that it is quite possible that such factorization will be a good approximation to reality.

The theory might be considered heuristic inasmuch as it is largely divorced from the more abstract theoretical developments and is based on a rather pragmatic extrapolation of experimentally observed phenomena. There are now many theories of high-energy interactions, and these are by no means all orthogonal. This model of high-energy processes is, perhaps, most similar to the "limited fragmentation" model of Yang and his colleagues. Indeed, it is possible that this diffraction-dissociation model can be taken under the wing of the Yang model.[8] The close relationship between the Yang model and the Regge model for high-energy diffractive processes has been emphasized by Hwa.[9]

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Comparison of Elastic $(\bar{p}p, pp)$, (K^-p, K^+p), and (π^-p, π^+p) Scatterings at ~8 and 16 GeV/c*

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Experimental data of $(\bar{p}p, pp)$, (K^-p, K^+p), and (π^-p, π^+p) elastic scatterings at ~8 and 16 GeV/c are examined in the context of the dual absorption model (DAM) as proposed by Harari. Results of this analysis of $(\bar{p}p, pp)$ and (K^-p, K^+p) exhibit similar behavior as expected from DAM and thus support the model. The behavior of (π^-p, π^+p), however, shows a marked difference from the other systems. An explanation for this difference is presented.

Well-known features of total and elastic cross sections of particles and antiparticles on protons above a few GeV/c can be generalized as follows:

(a) $\bar{\sigma}_{\text{tot}} > \sigma_{\text{tot}}, \quad \left.\frac{d\bar{\sigma}}{dt}\right|_{t=0} > \left.\frac{d\sigma}{dt}\right|_{t=0}$

and

(b) $\sigma_{\text{tot}}(s) \sim$ constant for exotic systems such as pp and K^+p.

Therefore, for a given s, crossover(s) of $d\bar{\sigma}/dt$ and $d\bar{\sigma}/dt$ for $t \neq 0$ can occur depending upon the behavior of $d\bar{\sigma}/dt$ and $d\sigma/dt$.[1] In this note, we present a detailed comparison of available experimental data of the elastic scattering of particles and antiparticles on protons, in particular, to explore the crossover phenomena in t in the framework of the dual absorption model (DAM) proposed by Harari.[2] In this model the imaginary part of the spin-nonflip, non-Pomeranchukon amplitude $R(t)$ can be extracted from elastic scattering of particles and antiparticles on protons as follows:

$$R(t) = \frac{\frac{d\sigma}{dt}(\bar{X}p \to \bar{X}p) - \frac{d\sigma}{dt}(Xp \to Xp)}{2\left(\frac{d\sigma}{dt}(Xp \to Xp)\right)^{1/2}}, \tag{1}$$

which is expected to have the general form

$$R(t) = Ae^{Bt}J_0(r\sqrt{-t}), \tag{2}$$

where $X = p$, K^+, and π^+ and J_0 is the Bessel function of zeroth order. Comparisons have been made by Davier and Harari[3] for 5-GeV/c $K^{\pm}p$ elastic data from counter and bubble-chamber experiments. The $R(t)$ has a zero at $t \sim -0.2$ GeV^2 (crossover in t), reaches a minimum at $t \sim -0.5$ GeV^2, and has another zero at $t \sim -1.3$ GeV^2. Thus, the behavior of $R(t)$ is quite similar to the J_0 function. In our investigation we examine the $R(t)$ function for the $(\bar{p}p, pp)$, (π^-p, π^+p) systems as well as (K^+p, K^-p) at ~8 and 16 GeV/c.

Ideally, the data used to evaluate $R(t)$ should all be taken with the same apparatus to minimize systematic errors, at the same value of beam momentum and with the same t binning for both par-