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MODELS OF FINAL STATES IN LEPTOPRODUCTION AND COLLIDING BEAMS*

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I. INTRODUCTION

I should first of all congratulate the organizers of this conference for their perspicacity in the choice of timing for this talk — immediately after Professor Richter's presentation[1] of the first SPEAR results — and in the choice of a title. As we shall shortly see, there is very little in the way of a real theory of final state hadrons in leptoproduction and colliding beams: there are only models.

But this is in part exactly what makes the subject so exciting and interesting. It is just here that rigorous light cone arguments don't apply, and more speculative additional assumptions are required to make predictions about the basic phenomena under discussion. For example, it is versions of the parton model which imply in some cases a "mixed scaling", i.e. scaling both in the Bjorken[2] variable ω associated with the presence of a current and in the Feynman[3] variable $x = p_L/p_L^{max}$, associated with a final hadron. Here one's intuition on what physics is fundamental and how to implement it is an important help — or hinderance.

A great many papers have been written in the subject area ostensibly covered by this talk. In the more than three preprint feet collected on my desk at the moment, one can find extensive applications of the vector dominance model, Regge-Mueller analysis, duality, light cone formalism, parton models, thermodynamic models, hydrodynamic models, etc., as well as combinations of these. Much of this is beyond the capacity of a simple country physicist like myself to understand. This is aside from the fact that in one hour it is simply impossible to summarize and make an adequate review of everything that has transpired in such a diffuse field.

Instead, I will try to give the general flavor of what has been done and to describe a few models of widely different viewpoint. Some approaches to the

subject as well as detailed amplification of the specific models presented are necessarily omitted. The talk has been left as it was before disclosure of the new e^+e^- data: there will be no instant analysis. I believe there is still much to be learned by examining the models already proposed in the light of what we know now.

II. VECTOR DOMINANCE MODEL

As a first manner of approach to the subject, let us consider the vector dominance model. I will particularly examine the work done by Greco et al.[4,5] because its comprehensive treatment of many different deep inelastic processes and its definiteness allows one to see more clearly the way the assumptions of the model fit together and interact with one another. Much related work has been done by Sakurai and collaborators.[6,7]

One starts by assuming an infinite number of vector mesons with a mass spectrum inspired by the Veneziano model,

$$m_n^2 = m_\rho^2(1 + an), \qquad n = 0, 1, 2, \ldots \tag{1}$$

where a is some constant. As we shall see below momentarily, photoproduction total cross sections demand that $a \simeq 2$. In this case, one has vector mesons at $m_\rho = 760$ MeV, $\sqrt{3}\, m_\rho \simeq 1300$ MeV, $\sqrt{5}\, m_\rho \simeq 1700$ MeV, and so forth. The state at $\sim \sqrt{3}\, m_\rho$ is the subject of a long search dating from the early days of the Veneziano model, while that at $\sim \sqrt{5}\, m_\rho$ might be identified with the observed ρ' (1600).

The process $e^+e^- \rightarrow$ hadrons is then pictured as proceeding by the virtual photon turning into a vector meson (coupling em_n^2/f_n for the n'th vector meson), which then decays into hadrons (total width Γ_n). The sum over all possible intermediate vector mesons gives the total amplitude. Being impressed by the

scaling observed in deep inelastic scattering, one forces scaling for $e^+e^- \to$ hadrons by demanding that

$$\Gamma_n \propto m_n \ , \tag{2}$$

and

$$m_n^2/f_n^2 = \text{constant} = b^2 = m_\rho^2/f_\rho^2 \ . \tag{3}$$

A direct calculation using Eqs. (2) and (3) and Breit-Wigner resonance forms for the vector meson intermediate states then produces the result[4,5]

$$\sigma_{e^+e^-}(Q^2) \propto \frac{1}{Q^2}\left(\frac{b^2}{a\,m_\rho^2}\right) = \frac{1}{Q^2}\left(\frac{1}{a\,f_\rho^2}\right), \tag{4}$$

where $Q^2 > 0$ is the (mass)^2 of the virtual photon.[8] Equation (4), by design, corresponds to the standard scaling behavior, $\sigma_{e^+e^-}(Q^2) \propto 1/Q^2$. Keeping track of 2's and π's yields with a = 2:

$$\sigma_{e^+e^-}(Q^2)/\sigma_{e^+e^- \to \mu^+\mu^-} = 8\pi^2/f_\rho^2 \sim 2.5 \ , \tag{5}$$

where $\sigma_{e^+e^- \to \mu^+\mu^-} = 4\pi\alpha^2/(3Q^2)$ is the (point) muon pair production cross section. A related, but more general, analysis by Sakurai[6] using a "new duality" principle obtains numbers between 3 and 5 for the right hand side of Eq. (5).

Photoproduction total cross sections are pictured in much the same way as those for annihilation: the photon becomes a vector meson (coupling em_n^2/f_n) which then strikes the nucleon (total cross section σ_n). One then again assumes a result motivated by scaling (in electroproduction):

$$\sigma_n \propto 1/m_n^2 \ . \tag{6}$$

The asymptotic photon-nucleon total cross section is then fixed and comes out as[4,5]

$$\sigma_{\gamma N} \propto \sigma_{\rho N}/(a^2 f_\rho^2) \quad . \tag{7}$$

It is here that a must be chosen as $\simeq 2$, for putting in $\sigma_{\rho N}$ as 28 mb then gives the correct value of $\sigma_{\gamma N}$.

All parameters in the model are now fixed. In particular, the structure functions in electroproduction are calculable in both shape and normalization. They scale,[2] i.e. are functions only of $\omega = 2M_N \nu/q^2$ as ν, $q^2 \to \infty$, as is to be expected from the enforcement of Eqs. (2), (3), and (6). As Fig. 1 shows, the result is fairly impressive, with $\nu W_2(\omega)$ having an asymptotic value (as $\omega \to \infty$) of 0.22.

This consistent little picture is completed by considering the single particle distributions expected in the photon fragmentation region in electroproduction "γ" + N $\to$ h(p') + ... As a result of all the previous scaling assumptions in Eqs. (2), (3), and (6), one finds a "super-scaled" form which only depends on dimensionless ratios,[4,5]

$$\frac{1}{\sigma} \frac{d\sigma}{(d^3p'/E')} \propto \left(\frac{1}{q^2}\right) G(\omega, x, p_\perp'^2/q^2) \ , \tag{8}$$

where $\omega = 2M_N \nu/q^2$ is the usual Bjorken scaling variable and $x = p_L'/p_L'^{max}$ is the Feynman variable representing the fractional longitudinal momentum carried by the produced hadron. The scaling in $p_\perp'^2/q^2$ results in

$$\langle p_\perp'^2 \rangle \propto q^2 \ , \tag{9}$$

a very strong prediction which has also been shown to follow in models of this sort using correspondence arguments by Bjorken.[9] An identical formula

(to Eq. (8)) has been obtained in the dual model by Gonzalez and Weis.[10] This differs importantly from the parton model results we shall discuss later, both in the scaling in $p'^2_\perp/q^2$ and in the factor of $1/q^2$ out in front of the right hand side. The results of the model have also been extended to such "mixed" processes as $pp \to \mu^+\mu^-$ + anything and various exclusive processes.[11,12]

Irrespective of the success or failure of the vector dominance model, it illustrates what is to me an important general method: One tries to see what constraints the regularities of hadrons force on the behavior of current induced processes, and conversely, what phenomena like scaling demand for relations among hadronic parameters like vector meson masses, couplings, and cross sections. A feature of positive value to some in such a picture is the ability to accommodate scaling in a fairly natural way without invoking point constituents.

A specific model like that outlined above, however, is rather tightly constrained. Predictions like $\langle p'^2_\perp \rangle \propto q^2$ are striking and very testable. While some increase in $\langle p^2_\perp \rangle$ with q^2 is probably observed[13] at low q^2 in the photon fragmentation region, it remains to be seen if such a dramatic growth in the hadrons' transverse momenta will continue to large q^2. Also, trouble in one place is liable to spread: recall the sensitivity to the parameter a, with $m_n^2 \propto a$, $\sigma_{\gamma N} \propto 1/a^2$, $\sigma_{e^+e^-} \propto 1/a$. For instance, lack of observation of a vector meson at $\sim\sqrt{3}\, m_\rho$ would change a and ruin the agreement in the value of $\sigma_{\gamma N}$ predicted. As new data becomes available it will be interesting to see how these are accommodated in such models.

III. STATISTICAL AND THERMODYNAMIC MODELS FOR ELECTRON-POSITRON ANNIHILATION

The early work of Bjorken and Brodsky[14] on a statistical model for e^+e^- annihilation has been improved and much extended, particularly in a recent

paper of Engels, Satz, and Schilling.[15] One may either start in the framework of the thermodynamic model of Hagedorn,[16] characterized by a maximum hadron temperature T_0, or work in terms of the statistical bootstrap[17] applied to cascade decays from one hadronic "fireball" to another, with asymptotically bounded average energy of a secondary hadron. Both these approaches lead to the same result asymptotically.[15]

A principal characteristic of such models for e^+e^- annihilation is that the average final hadron's energy, $<E'>$, approaches a finite limit as $Q^2 \to \infty$. For a hadron temperature $kT_0 = 160$ MeV (from strong interaction production processes), one finds $<E'> = 320$ MeV for production of massless hadrons and $<E'> = 414$ MeV for real pions. The finite value of the average energy per produced particle means that

$$<n> \propto \sqrt{Q^2} , \tag{10}$$

a growth with the highest power of Q^2 allowed by energy conservation.

Single particle distributions have the simple form $e^{-cE'}$ where c is a predicted constant which is independent of Q^2. An example of the finite Q^2 predictions, which have been worked out in detail,[15] is shown in Fig. 2. Such an exponential behavior near the kinematic boundary[18] ($E' \simeq \sqrt{Q^2}/2$) is in disrepute with some theorists, as it would imply a similar exponential behavior of form factors from duality or correspondence arguments.[19]

Leaving aside questioning the basic theoretical assumptions of such models, they are difficult to prove or disprove experimentally with a limited Q^2 range on the basis of the behavior of $<E>$ or $<n>$. In this regard the situation is similar to that of a few years ago with respect to the diffractive excitation or nova model in purely hadronic collisions.[20] There the decisive test[21] came with two particle

correlations; such measurements will very likely also be crucial here. We will return to this question again at the end of the talk.

IV. THE HYDRODYNAMIC MODEL

A somewhat related theoretical approach to e^+e^- annihilation is given by the Landau hydrodynamic model, which also has previously been applied to strong interactions and electroproduction.[22] The basic idea is that the energy involved in a collision is dumped into a localized region, V_0, where there are many field quanta present and a corresponding small mean free path for interaction. As a result one has a hot gas or fluid, which expands according to the laws of hydrodynamics. As it expands, it cools, and the energy density drops until it reaches $\sim m_\pi c^2$ in a pion Compton volume, at which point individual hadrons (pions) condense out to form the observed final state.

Some manipulation of thermodynamic equations shows that[23] in this situation the number of particles produced, $<n>$, goes like

$$<n> \propto V_0^{1/4} (\sqrt{s})^{3/4} , \tag{11}$$

where $\sqrt{s}$ is the total center of mass energy involved. For electroproduction, one works in the virtual photon-nucleon center of mass and, given the Lorentz contraction along the direction of motion, assumes that the initial volume $V_0 \propto M_N/E_N$. This results in

$$<n> \propto \frac{s^{1/2}}{(s + M_N^2 + q^2)^{1/4}} . \tag{12}$$

For fixed q^2 and large s one recovers the familiar $s^{1/4}$ behavior of $<n>$, as for hadron-hadron collisions in the hydrodynamic model.[22] Note though that at fixed s, $<n>$ decreases as $(1/q^2)^{1/4}$ as $q^2 \to \infty$.

For the case of e^+e^- annihilation into hadrons, V_0 is assumed[22, 24, 25] to be constant with $s = Q^2$, so that

$$\langle n \rangle \propto (\sqrt{Q^2})^{3/4} = (Q^2)^{3/8} . \tag{13}$$

Correspondingly, there is a slow rise with Q^2 of the mean energy per hadron:

$$\langle E' \rangle \propto (Q^2)^{1/8} , \tag{14}$$

Detailed calculations of particle distributions have been done by Cooper et al.,[24] who find the single particle distribution to be that of a Lorentz boosted (from the hydrodynamic expansion) Bose distribution characterized by $kT = m_\pi c^2$ in its rest frame. The value of $\langle E' \rangle$ determines the boost, with there being no boost at all for $\langle E' \rangle \simeq 430$ MeV. Present data[1] on $\langle E' \rangle$ indicates a very slight boost, and correspondingly small hydrodynamic expansion as is seen in Fig. 3.[24]

V. SOME GENERAL CONSIDERATIONS

Let us digress for a moment from our considerations of particular models to discuss some more general ideas which should apply to all models. One such principle is the inclusive-exclusive connection. While originally used to derive the behavior of inclusive hadronic amplitudes near the phase space boundary,[3] it was first applied to inelastic electron scattering by Bloom and Gilman[26] where the principle may be realized in a duality framework. It has been generalized and applied in many other situations by Bjorken and Kogut.[27]

The idea is most simply illustrated in deep inelastic scattering[26] by considering the behavior of the inclusive structure function νW_2 near $\omega = 1$. Assuming a power law in $\omega - 1$,

$$\nu W_2(\omega) \propto (\omega - 1)^p , \tag{15}$$

we note that this can be rewritten in terms of the final hadronic invariant mass, W, and q^2 as

$$\nu W_2^{\text{inclusive}} \propto (W^2/q^2)^p . \tag{16}$$

Now focus instead on a particular resonance (or more generally), mass interval) which has a form factor $G(q^2)$ which behaves as

$$G(q^2) \rightarrow (1/q^2)^{n/2} \tag{17}$$

as $q^2 \rightarrow \infty$. Then at large q^2 it makes a contribution to νW_2 which is

$$\nu W_2^{\text{exclusive}} \propto q^2 |G(q^2)|^2 \rightarrow (1/q^2)^{n-1} . \tag{18}$$

Matching up powers of q^2 between (16) and (18) immediately gives

$$n = p + 1 , \tag{19}$$

which is the familiar result of Drell and Yan[28] and West[29] for the elastic form factor and of Bloom and Gilman[26] for excitation form factors of nucleon resonances in general. Conversely, if we assume the general validity of Eq. (19), we have a way of getting at the power law involved in elastic or resonance form factors by measuring the shape of the inclusive spectrum.

Another interesting general idea is that of "crossing" from the space-like to time-like region near threshold. While for certain class of graphs the structure functions in the two regions are analytic continuations of one another, as suggested by Drell, Levy, and Yan,[30] in general this is not true due to certain "double discontinuity" terms.[31,32,33] However, for a restricted region near threshold, a connection[30] between the two regions may still exist. Recently Gatto et al.[34] have made such a connection very plausible by examining large classes of graphs in field theory.

To see what has been shown, consider the variable $\omega = -2p \cdot q/q^2$, which has the domain $0 \leq \omega \leq 1$ for $e^+e^- \to \bar{h}(-p) + \ldots$, and $1 \leq \omega \leq \infty$ in $e + h(p) \to e + \ldots$. We are interested in the point $\omega = 1$ where the two domains touch. If $\pm\bar{F}(\omega) \to c_1(1 - \omega)^{p_1}$ as $\omega \to 1^-$ in e^+e^- annihilation[35] and $F(\omega) \to c_2(\omega - 1)^{p_2}$ as $\omega \to 1^+$ in electroproduction, then the result of Gatto et al.,[34] is that $p_1 = p_2$ and $c_1 = c_2$. Thus with restricted forms for the structure functions near $\omega = 1$ one can continue from the space-like to time-like region in the immediate neighborhood of $\omega = 1$, but not generally.[36]

This leads to the very interesting possibility of measuring the asymptotic behavior of any hadron's form factor, $F_h(q^2)$, using e^+e^- annihilation: the shape of $e^+e^- \to h + \ldots$ near $\omega = 1^-$ could be continued to give the shape of $eh \to e + \ldots$ near $\omega = 1$ and thence the inclusive-exclusive connection gives the power law behavior of $F_h(q^2)$ as $q^2 \to \infty$. As better e^+e^- data becomes available this could be of great importance in understanding hadron structure, particularly the relation between that of mesons and that of baryons.

VI. QUARK PARTON MODEL

Next we turn to the quark parton model as developed by Feynman,[3] Bjorken,[37] and others. There are two basic reasons for considering quarks as the partons in my view. First is the success of the quark model for hadrons at low energy. Not only does the spectrum of states show easily identifiable SU(6) × O(3) multiplets, but recent developments[38] indicate that photon and pion decay amplitudes have a structure such as has been abstracted from the quark model.[39] In particular, there are more than 25 relative signs of amplitudes for $\gamma N \to N^* \to \pi N$ and $\pi N \to N^* \to \pi\Delta$ which are correctly predicted. Second, quark partons provide an exceedingly simple description of both eN and νN deep inelastic scattering. If I may be allowed to greatly simplify and condense some of the data[40] we heard

yesterday, then not only does deep inelastic scattering behave as if the nucleon was composed of point, spin 1/2 objects with quantum numbers consistent with those of quarks, but there is at most a very small anti-quark component.[41] That is not to say that at sufficiently high ω we will not see a sizeable $q\bar{q}$ component (corresponding to Pomeron exchange), but in the presently well explored range of ω, say $1 < \omega < 10$ the data (particularly $\sigma_T^{\bar{\nu}N}/\sigma_T^{\nu N}$) show that the nucleon acts as if composed of just three quarks.

Consider then the possible rapidity regions available for a final state hadron, h(p'). We define the rapidity as in hadronic reactions,

$$y = \frac{1}{2} \ell n \frac{E' + p'_L}{E' - p'_L} , \tag{20}$$

where p'_L is the momentum along the beam direction. For a process like $pp \to \pi +$ anything we expect (at fixed $p'_\perp$) on the basis of a model with short range correlations that the invariant cross section $E' \, d\sigma/d^3p'$ will exhibit a behavior like that in Fig. 4a. At a given value of $p'_\perp$ (which is limited) there are fragmentation regions of fixed finite length in rapidity associated with the beam or target, separated by a central plateau of constant height. As the total rapidity interval available is of length $\ell n \, s$ and the fragmentation regions are of fixed finite length, we see that the central plateau is $\ell n \, s$ in length. The multiplicity, which is proportional to the area under the curve, then grows as the length of the central plateau grows, i.e. as $\ell n \, s$.

An analogous result is expected for the photoproduction process $\gamma p \to \pi + \ldots$, as shown in Fig. 4b. In fact, normalizing to the total cross section, the proton fragmentation region and central plateau should be just as before. Only the beam fragmentation region in Fig. 4b (of fixed, finite length in rapidity) has changed from that in Fig. 4a to one characteristic of the photon.

Now what happens when q^2 varies? The situation within the Regge-Mueller[42] framework is shown[43] in Fig. 5. The nucleon fragmentation region is as before, while the central plateau has the same height, but its length is now $\ell n\, \omega$. The remaining $\approx \ell n\, s - \ell n\, \omega \approx \ell n\, q^2$ of available rapidity interval comes from photon (or more generally, current) fragmentation. In Fig. 5 we have three subdivisions: the parton and hole fragmentation regions of finite length, and a current plateau of length $\sim \ell n\, q^2$. While the existence of the regions is general,[44,45] the way they are drawn and the names in Fig. 5 are specific to the parton model. There the hole fragmentation region[46] is the location in rapidity of the struck parton before the interaction, while the parton fragmentation region is its location afterward.[46,47] The current plateau may or may not[48] exist, and its height may or may not be equal to that of the central plateau. That there is a region of rapidity of length $\ell n\, q^2$ associated with the current is a general kinematic fact however. Moreover, in e^+e^- colliding beams one would similarly expect a current plateau of length $\ell n\, q^2$, with parton fragmentation regions of finite length at the ends. The question now before us is the specific content of these regions and their interrelation in the quark parton model.

In the parton fragmentation region the quark parton model is implemented with the following assumptions:[47,49]

1. The virtual photon interacts with point, spin 1/2 constituents, the quark partons. This is the standard assumption which yields scaling of νW_2 and W_1.

2. The parton fragments into hadrons, independently of how it is produced. This will relate final state hadrons in eN, νN, and e^+e^- collisions.

3. The hadron distribution from a given parton is only a function of $z = p^{(\text{hadron})}/p^{(\text{parton})}$ and $p_\perp$ of the hadron relative to the parton direction. The $p_\perp$ distribution is assumed to be limited, so that at high enough energies there will be jets along the parton direction of motion on an event by event basis.

Summarizing the parton (of type i) fragmentation into a given hadron (h) by the function $D_i^h(z)$, the final hadron distribution in $e + N \rightarrow e + h + \ldots$ after integrating out the transverse momentum dependence is given by[50]

$$\nu \, d\sigma_T/(dp/E) \propto \sum_i Q_i^2 \, f_i(1/\omega) \, D_i^h(z) \, . \tag{21}$$

Here Q_i is the charge and $f_i(1/\omega)$ is the momentum distribution of the i'th parton in the nucleon, so that

$$W_1 \propto \nu \sigma_T \propto \sum_i Q_i^2 \, f_i(1/\omega) \tag{22}$$

is the standard result[50] for the structure function W_1. Equation (21) shows the "mixed" scaling characteristic of the parton model predictions: the inclusive distribution depends only on the Bjorken variable ω and on z, which is essentially the usual Feynman variable[3] x in this case.

Similarly, for e^+e^- annihilation into hadrons, where the parton model predicts

$$\sigma_T \propto \left(\sum_i Q_i^2\right) \Big/ Q^2 \, , \tag{23}$$

the distribution of hadron h should be

$$d\sigma/(dp/E) \propto \sum_i Q_i^2 \, D_i^h(z)/Q^2 \tag{24}$$

where now $z = |\vec{p}|/(\sqrt{Q^2}/2)$.

There are clearly close connections between the parton fragmentation regions[49] of eN, νN, and e^+e^- scattering since the same functions $D_i^h(z)$ appear in Eqs. (21) and (24). This is made even more clear by noting that there are

relations among the $D_i^h(z)$'s following from isospin and charge conjugation invariance. For example, labeling the quarks as u, d, and s, one has

$$D_u^{\pi^+}(z) = D_d^{\pi^-}(z) = D_{\bar{u}}^{\pi^-}(z) = D_{\bar{d}}^{\pi^+}(z) , \tag{25}$$

and

$$\frac{1}{2}\left[D_i^{\pi^+}(z) + D_i^{\pi^-}(z)\right] = D_i^{\pi^0}(z) , \tag{26}$$

where i = u, d, s, $\bar{u}$, $\bar{d}$, or $\bar{s}$. This last result follows from the quarks having isospin $I \leq 1/2$. Equations (25) and (26) and their analogues lead to sum rules for the final hadrons observed in eN and νN inelastic scattering.[49] Equation (26) demands that

$$\frac{1}{2}\left[d\sigma^{\pi^-}(z) + d\sigma^{\pi^+}(z)\right] = d\sigma^{\pi^0}(z) , \tag{27}$$

in the parton fragmentation region for each value of z. Applied to e^+e^- annihilation, where the distribution of π^+'s and π^-'s must be equal by charge conjugation invariance, one obtains

$$d\sigma^{\pi^-}(z) = d\sigma^{\pi^+}(z) = d\sigma^{\pi^0}(z) . \tag{28}$$

While very high energies and q^2 values are required to really see the parton fragmentation region clearly, such a model provides a natural "explanation" for the new qualitative features of hadron electroproduction at SLAC energies. The most striking such feature is the increase in the π^+/π^- ratio for pions produced on protons in the virtual photon's direction with a finite fraction of its momentum[13,51] (which corresponds at high energies to the parton fragmentation region), as shown in Fig. 6. In a quark-parton picture this comes about because the u quarks are much more likely than d quarks to be struck by the virtual photon and thrown

forward, where they preferentially fragment into π^+'s rather than π^-'s. A number of detailed fits[52] to the data have been made using the single arm electroproduction data to obtain the $f_i(1/\omega)$, and the hadron electroproduction data to then fit the $D_i^h(z)$. The qualitative points made above are borne out, and specific predictions for other processes (νN, e^+e^-) are obtained as well.[53]

An interesting sidelight to the increase of the π^+/π^- ratio with q^2 in electroproduction is that a similar effect[51] has been seen (Fig. 7) in the photon fragmentation region with increasing $p_\perp$ in photoproduction ($q^2 = 0$). Very recent measurements[54,55] show that the ratio $\pi^o/\frac{1}{2}(\pi^+ + \pi^-)$ also <u>increases</u> with $p_\perp$ and is greater than unity for $p_\perp^2 \sim 1\ \mathrm{GeV}^2$. It is clearly of great interest to see what happens then in electroproduction. If the ratio $\pi^o/\frac{1}{2}(\pi^+ + \pi^-)$ also exceeds unity as q^2 increases, then Eq. (27) will be violated and we have been fooling ourselves with quark-parton explanations of the π^+/π^- ratio. The fact that the neutral energy considerably exceeds half the charged particle (mostly pions) energy at SPEAR[1] is a very troubling sign in this regard. An experiment[56] to measure inclusive π^o electroproduction is now underway at SLAC.

Let us now turn to the current plateau region of Fig. 5. Without such a plateau, quark partons would be produced in isolation in the parton fragmentation region. Indeed, exactly such an isolation occurs in the multiperipheral model,[57] softened field theories,[58] and early versions of the covariant parton model.[59,60]

The possibility of cascade decay of a parton to produce such a plateau has been investigated in detail using space-time arguments by Kogut, Sinclair, and Susskind.[61] They find that such a mechanism, of the so-called "outside-inside" type, is not sufficient to generate a plateau and prevent isolated quark quantum numbers from appearing.[62] Instead recent work has proposed an "inside-outside"

mechanism.[9] In deep inelastic eN and νN scattering, $q\bar{q}$ pairs, produced by the vacuum polarization between the parton and the hole, move outward, filling up the current plateau until they catch up with the struck parton and neutralize its quantum numbers.[63] An identical mechanism applies to e^+e^- annihilation where the generation of the plateau starts at the center and catches up to the parton in a time $\sim\sqrt{Q^2}$.

Such a picture is explicitly realized[64] in quantum electrodynamics in two dimensional space-time, where there is total shielding of the fermion's charge (analogue of quark quantum numbers ?), although the short distance properties of the theory are those of free fermions. The jump to four dimensions with gauge vector bosons coupled to the quark quantum numbers which one wants totally screened (e.g. triality) is obvious. The electromagnetic vacuum polarization in such a world is pictured by Bjorken[9] as in Fig. 8. It is here that recent work on "asymptotic freedom" in gauge theories may be of particular relevance, producing (almost) scaling behavior in a field theory context.[65] Much work along these lines is in progress.[66]

Accepting for the moment the existence of a mechanism to generate such a plateau, we note that its height should be universal[9,47] in eN, νN, and e^+e^- annihilation. Thus the mean multiplicity in annihilation should be

$$\langle n \rangle_{e^+e^-} = C_{e^+e^-} \ln Q^2 + \text{const.} , \tag{29}$$

assuming a constant height of the plateau and fragmentation regions of finite length. The analogous expression in deep inelastic scattering is[67]

$$\langle n \rangle_{eN} = C_{e^+e^-} \ln q^2 + C_h \ln \omega + \text{const.}, \tag{30}$$

where C_h is the height of the ordinary hadron (central) plateau.

An important question arises at this point of whether the current and central plateau are of the same nature. Cahn and Colglazier[68] have argued that they may well be different: the current plateau has quarks on the ends which lead to a different polarization of the intervening medium than for the hadron case. If the plateaus have the same height, $C_{e^+e^-} = C_h$, then Eq. (30) collapses to

$$\begin{aligned} \langle n \rangle_{eN} &= C_h \ln q^2 \omega + \text{const.} \\ &\simeq C_h \ln s + \text{const.}, \end{aligned} \tag{31}$$

just as for real photons or hadrons. Recent Cornell data,[69] although at very low s values for application of such ideas, supports Eq. (31).

Most of this picture is realized in the work of Preparata et al.[70] on the massive quark model. The scaling behavior in eN, νN, and e^+e^- annihilation, as well as the various regions of rapidity (including the current plateau) are all explicit. In particular, the current and central plateaus have the same composition but different heights: $C_{e^+e^-} \leq C_h$. In addition, in several cases one can calculate the approach to the various scaling limits using exchanges with $\alpha(0) < 1$; Pomeron exchange (with $\alpha(0) = 1$) generates the scaling behavior itself.

VII. WHAT CAN WE LOOK FOR IN THE IMMEDIATE FUTURE

While it would be nice to be able to look at the whole structure of the fragmentation and plateau regions of the "limosine" in Fig. 5, we are a long way from being able to do so. Judging from NAL and ISR data, which suggest that something like six units of rapidity are necessary to study such structure with a decent separation of the various regions in pp collisions, we see that for current induced processes we should like $\ln \omega \simeq 6$ and $\ln q^2 \simeq 6$ or $\ln s \simeq 12$, i.e.,

$s \simeq 160,000\ GeV^2$! The effect of the energy crisis in which we find ourselves on Fig. 5 is shown[71] in Fig. 9.

But there are many things which are presently accessible. The question of scaling in the parton fragmentation region is accessible to a good test, and while there are hints that forward going pions do scale in the required way[51] (Eq. (21)) for ep inelastic scattering, much remains to be done. In particular, the relations we have discussed between the particle composition and distribution in eN, νN and e^+e^- deep inelastic processes remain basically unchecked.

Then there is the question of the plateau in e^+e^- annihilation. The generation of storage rings already built or under construction should give us a good indication of its properties, particularly its height, if it exists at all. The question of the existence of jets should also be accessible to such e^+e^- experiments.

Also testable in the immediate future are the ratios of various particles in the fragmentation regions of e^+e^-, eN, and νN, versus similar hadronic ratios. In particular we saw above that the ratio of $\pi^+/\pi^-/\pi^o$ may be a first indication of trouble for the quark parton model. In general, investigation of such ratios should tell us much about the basic mechanism underlying the production of the final state hadrons.

One of the important ways of distinguishing among different mechanisms or models, just as it is for hadronic reactions, is investigation of two particle correlations. There has been a considerable amount of work on such questions in the various models I have discussed.[72] As an example, consider the average charge transfer squared, $\langle u^2 \rangle$, between hemispheres in e^+e^- annihilation. As discussed by Newmeyer and Sivers,[73] a quark parton model with hadrons emitted from quarks forming a closed loop results in small local charge fluctuations and a bounded $\langle u^2 \rangle$. A statistical model, on the other hand, has large fluctuations

which grow as $<n>$. This serves to illustrate in another way the difference between careful planning and fireballs. At higher energies the difference between the predicted values of $<u^2>$ in the two models is large and testable, as shown in Fig. 10.

In preparing this talk I happened to look at the proceedings of the previous conference in this series.[74] At the end of that conference there was a panel discussion, the transcription of which still seems very much worth reading. At one point, in reply to a challenging question of J.D. Jackson, it was noted by Feynman that the most important advances in our understanding often come out of apparent paradoxes which occur when we are presented with two sets of facts which seem mutually impossible.

Looking back at some of the models I have discussed, one also notes our tendency to follow the same old paths unless shaken out of them. In this regard, I am also reminded of some remarks attributed[75] to Sidney Coleman in an interview by a science fiction magazine. When asked about the relative use of imagination in science fiction versus that in real science, he replied to the effect that there was more use of imagination in science, as one is constrained by the facts to be imaginative. It certainly appears that deep inelastic scattering and annihilation into hadrons are giving us such facts. Perhaps they are even the paradoxical facts for which Feynman is looking. Now it is time to go and think about the latest exciting and very puzzling data presented at this conference.

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FIGURE CAPTIONS

1. The structure function νW_2 calculated from the vector dominance model of Refs.4 and 5. The quantities νW_2^D and νW_2^R are the diffractive and resonant contributions to νW_2, respectively (see Ref. 5).
2. Single particle distributions for e^+e^- annihilation at $\sqrt{Q^2} = 6$ GeV in the thermodynamic model of Ref. 15.
3. Single particle distribution of pions in e^+e^- annihilation at $\sqrt{Q^2} = 3$ and 5 GeV in the hydrodynamic model of Ref. 24.
4. Schematic rapidity distributions of pions produced in the reactions: (a) $pp \rightarrow \pi$ + anything, and (b) $\gamma p \rightarrow \pi$ + anything.
5. Schematic rapidity distribution of particles produced in virtual photon-nucleon collisions.
6. Ratio of positive to negative hadrons produced[51] with $0.4 < x < 0.85$ on proton and neutron targets as a function of q^2.
7. Ratio of positive to negative pions produced[51] by real photons incident on protons as a function of $p_\perp^2$.
8. Vacuum polarization in a theory[9] with gauge vector mesons (wiggly lines) and quarks (solid lines) whose quantum numbers are totally shielded by the "backflowing" inner quark current.
9. A possible rapidity distribution at energies available in the forseeable future.
10. The average charge transfer squared, $\langle u^2 \rangle$, between hemispheres in e^+e^- annihilation in various models.[73]

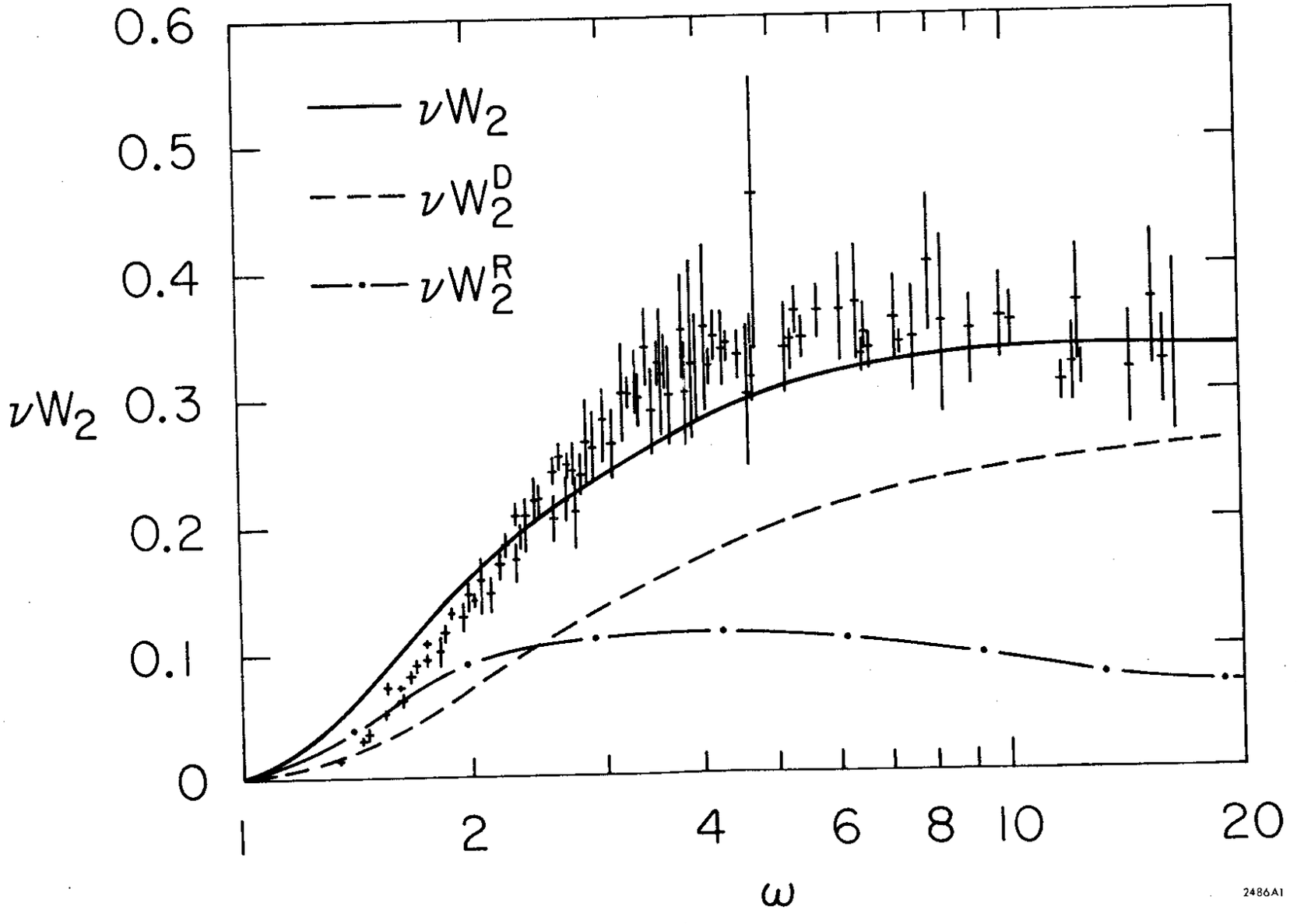
0.6
0.5
0.4
0.3
0.2
0.1
0
νW_2
νW_2
νW_2^D
νW_2^R
1
2
4
6
8
10
20
ω
2486A1

Fig. 1

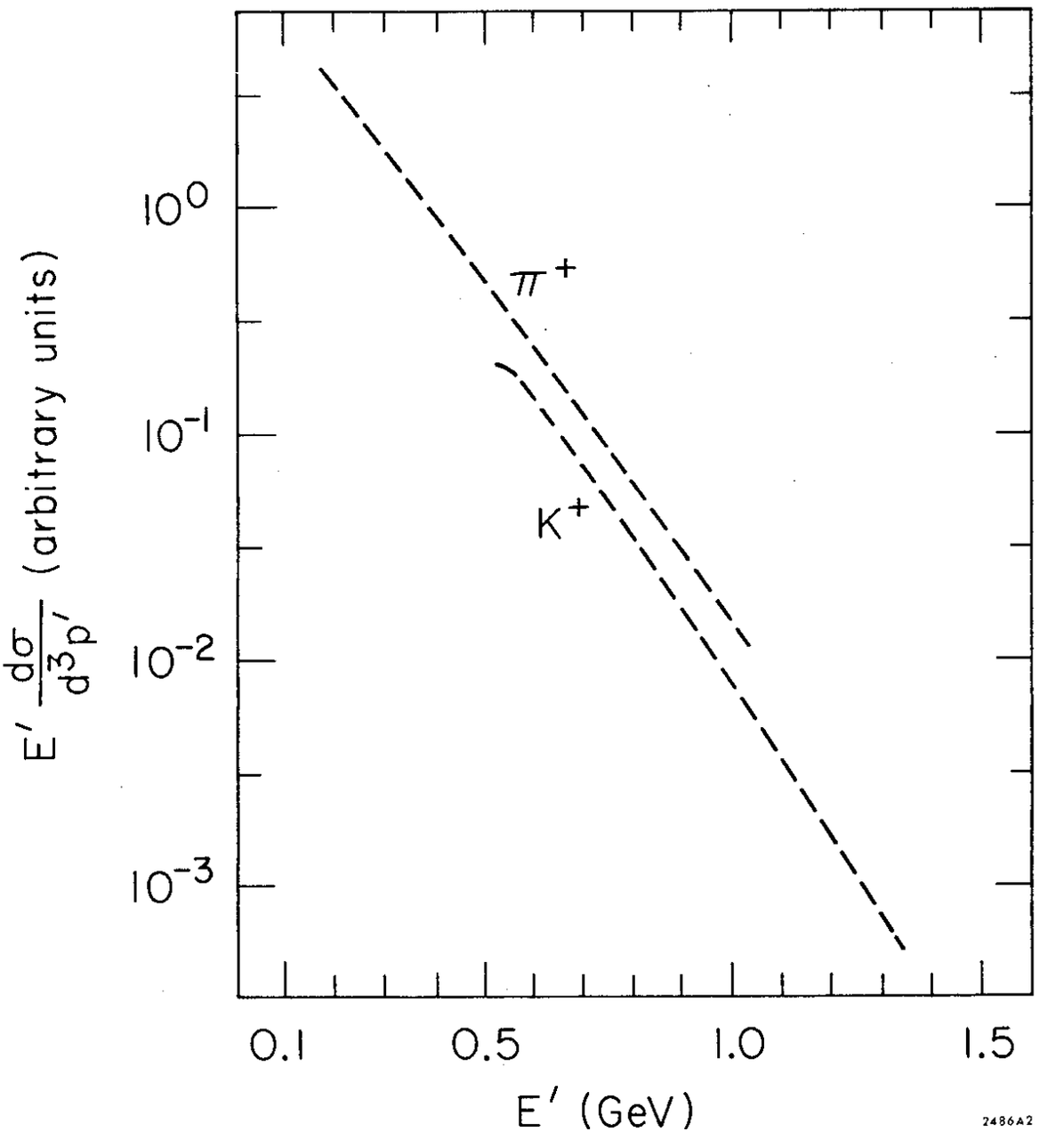
$E' \frac{d\sigma}{d^3p'}$ (arbitrary units)
10^0
10^{-1}
10^{-2}
10^{-3}
π^+
K^+
0.1
0.5
1.0
1.5
E′ (GeV)
2486A2

Fig. 2

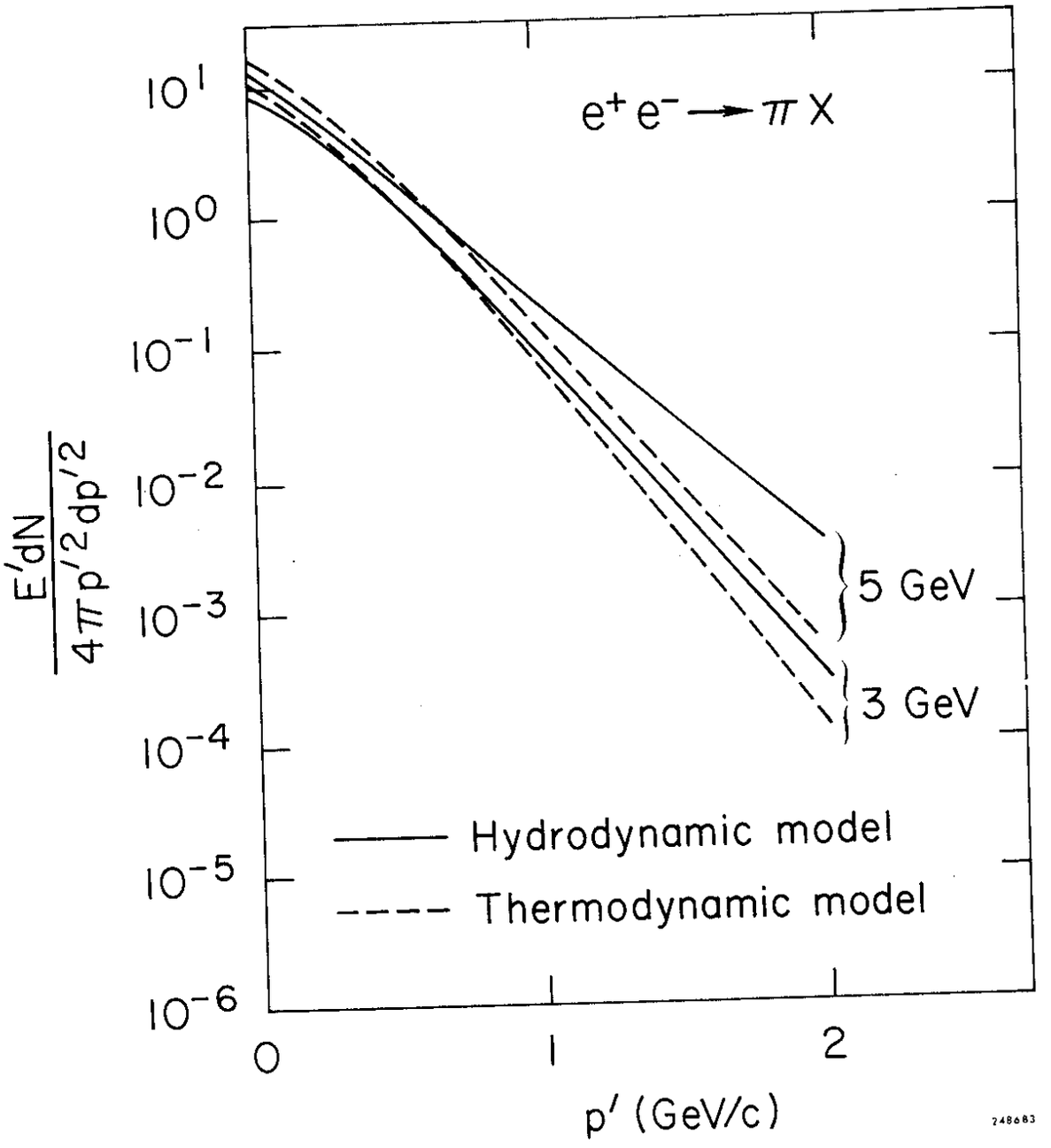
$e^+ e^- \rightarrow \pi X$
10^1
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
$\frac{E' dN}{4\pi p'^2 dp'^2}$
5 GeV
3 GeV
Hydrodynamic model
Thermodynamic model
0
1
2
p' (GeV/c)
248683

Fig. 3

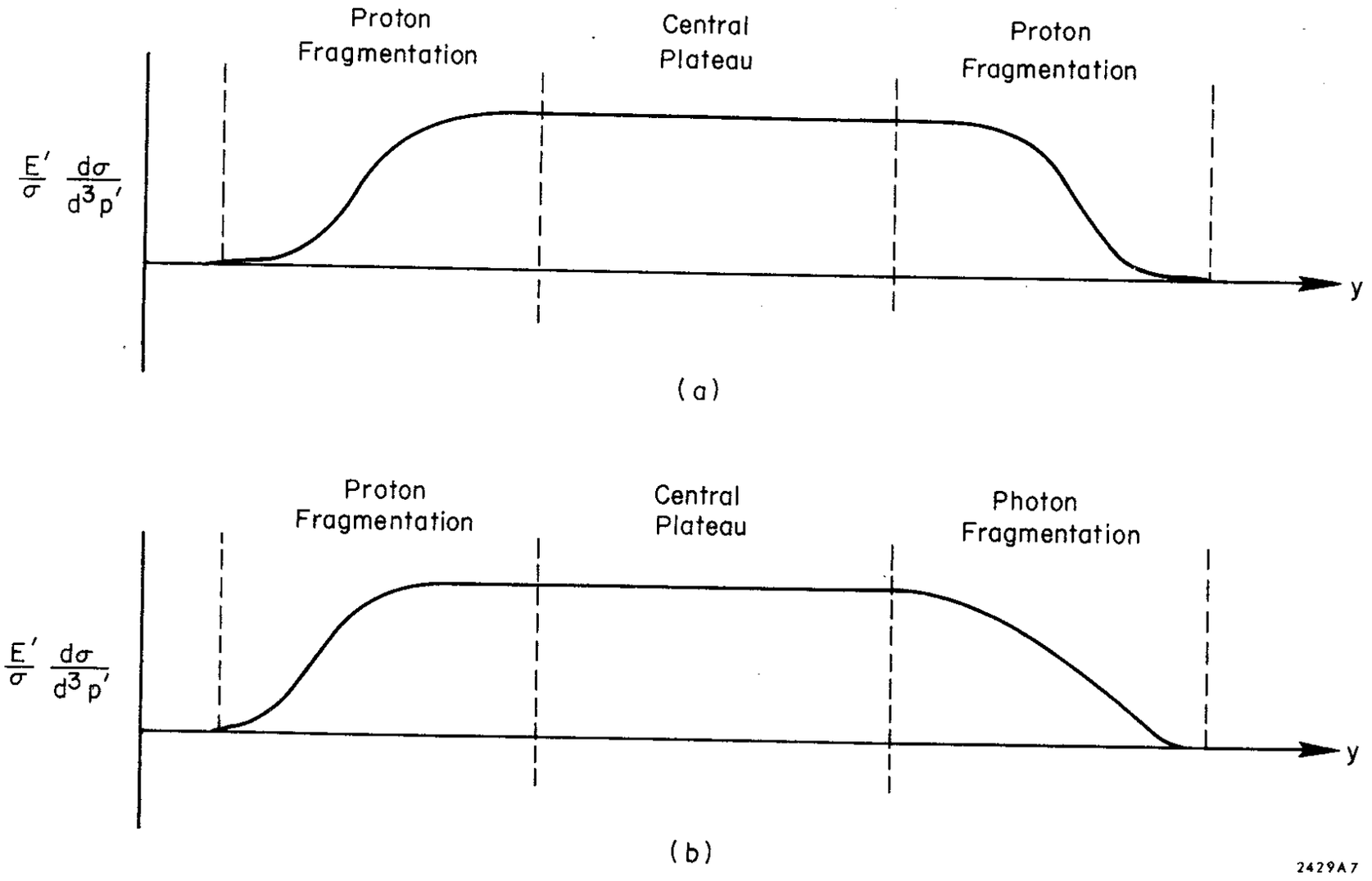
Proton
Fragmentation
Central
Plateau
Proton
Fragmentation
$\frac{E'}{\sigma}\frac{d\sigma}{d^3p'}$
y
(a)
Proton
Fragmentation
Central
Plateau
Photon
Fragmentation
$\frac{E'}{\sigma}\frac{d\sigma}{d^3p'}$
y
(b)
2429A7

Fig. 4

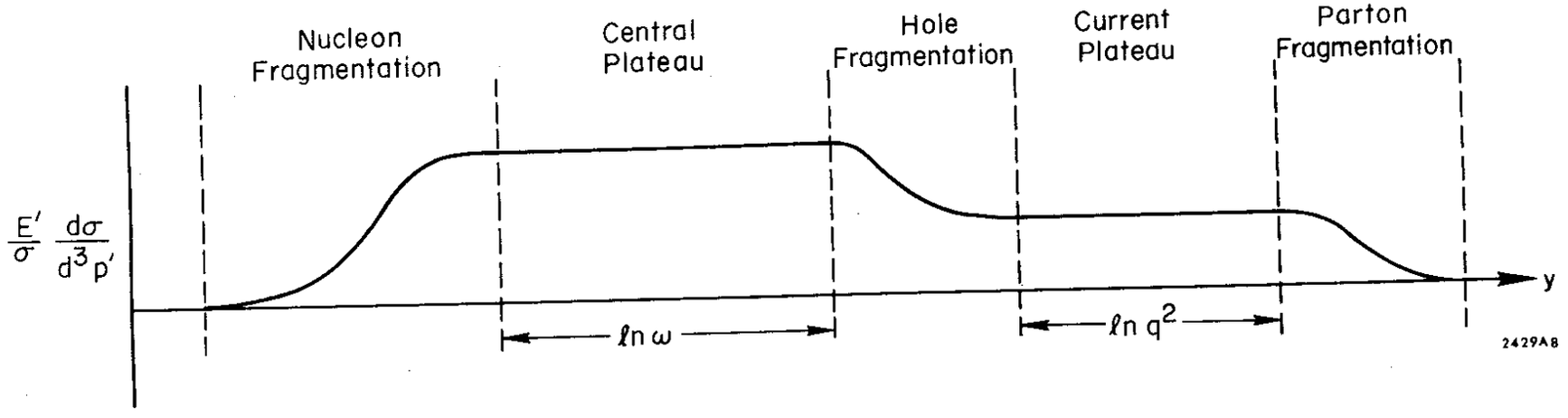
Nucleon Fragmentation
Central Plateau
Hole Fragmentation
Current Plateau
Parton Fragmentation
$\frac{E'}{\sigma}\frac{d\sigma}{d^3p'}$
y
$\ell n\, \omega$
$\ell n\, q^2$
2429A8

Fig. 5

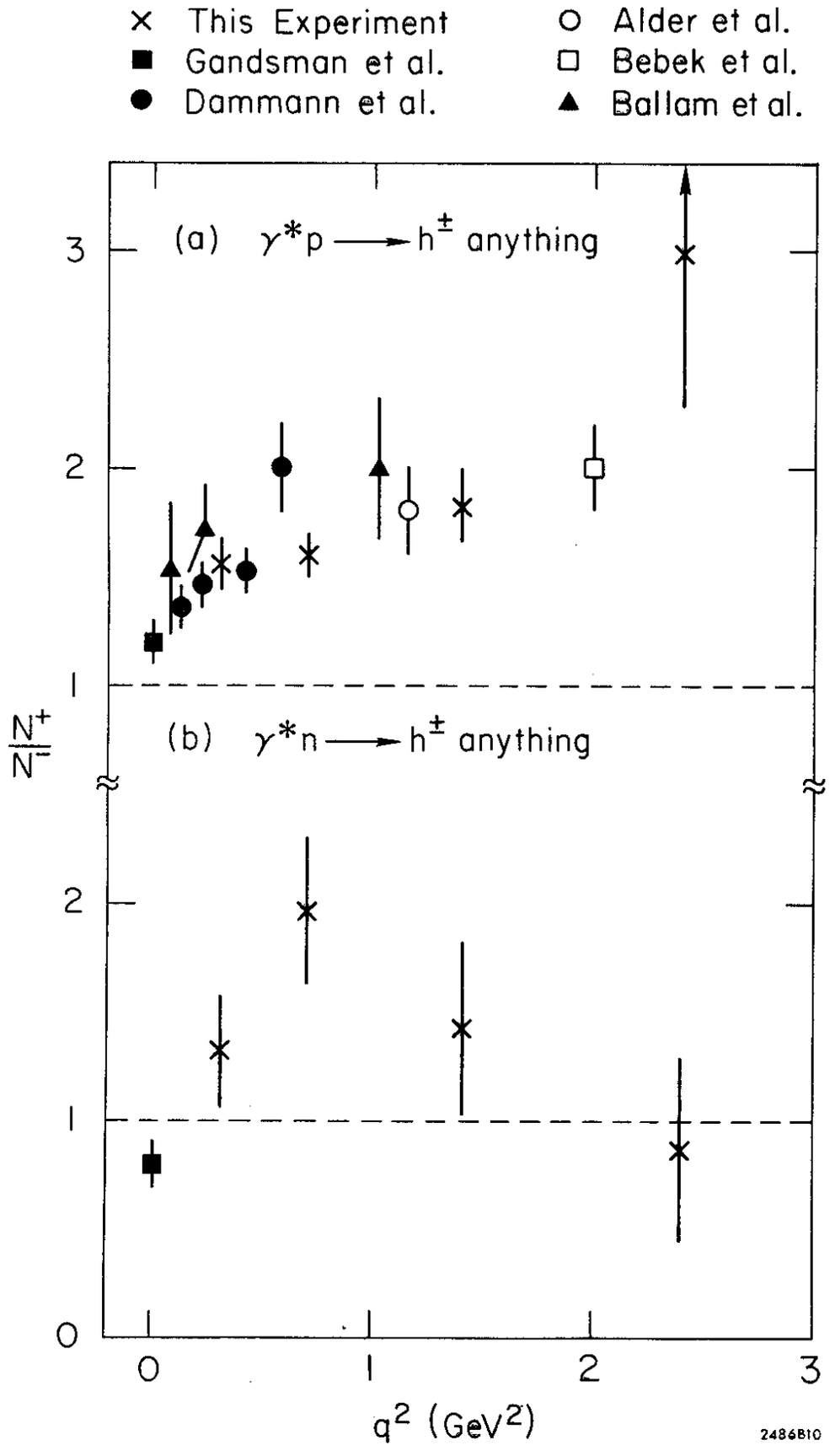
× This Experiment
■ Gandsman et al.
● Dammann et al.
○ Alder et al.
□ Bebek et al.
▲ Ballam et al.
(a) $\gamma^* p \longrightarrow h^{\pm}$ anything
(b) $\gamma^* n \longrightarrow h^{\pm}$ anything
$\frac{N^+}{N^-}$
q^2 (GeV^2)
2486B10

Fig. 6

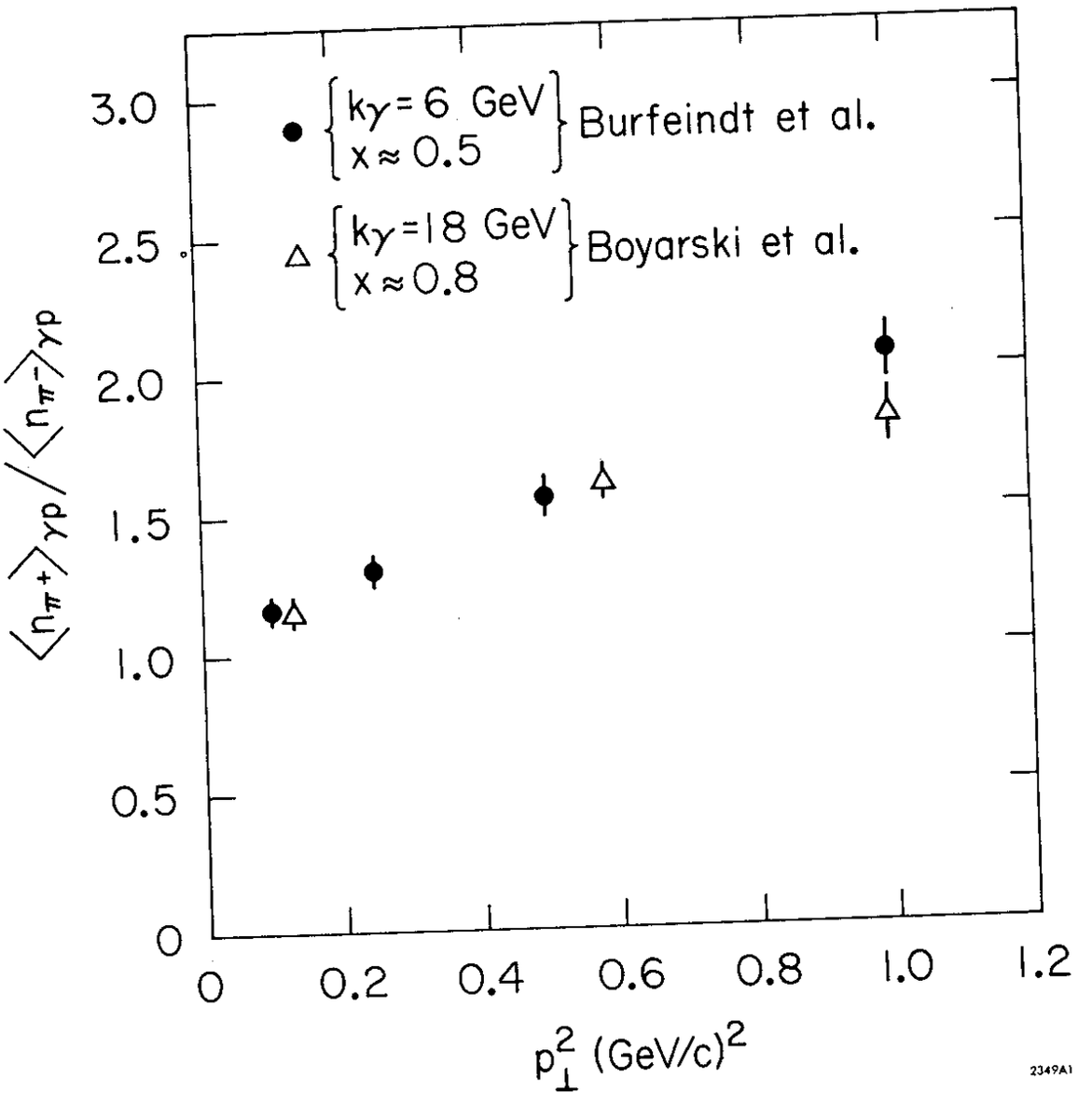
3.0
2.5
2.0
1.5
1.0
0.5
0
0
0.2
0.4
0.6
0.8
1.0
1.2
$\langle n_{\pi^+} \rangle_{\gamma p} / \langle n_{\pi^-} \rangle_{\gamma p}$
$p_\perp^2$ (GeV/c)^2
$k\gamma = 6$ GeV
$x \approx 0.5$
Burfeindt et al.
$k\gamma = 18$ GeV
$x \approx 0.8$
Boyarski et al.
2349A1

Fig. 7

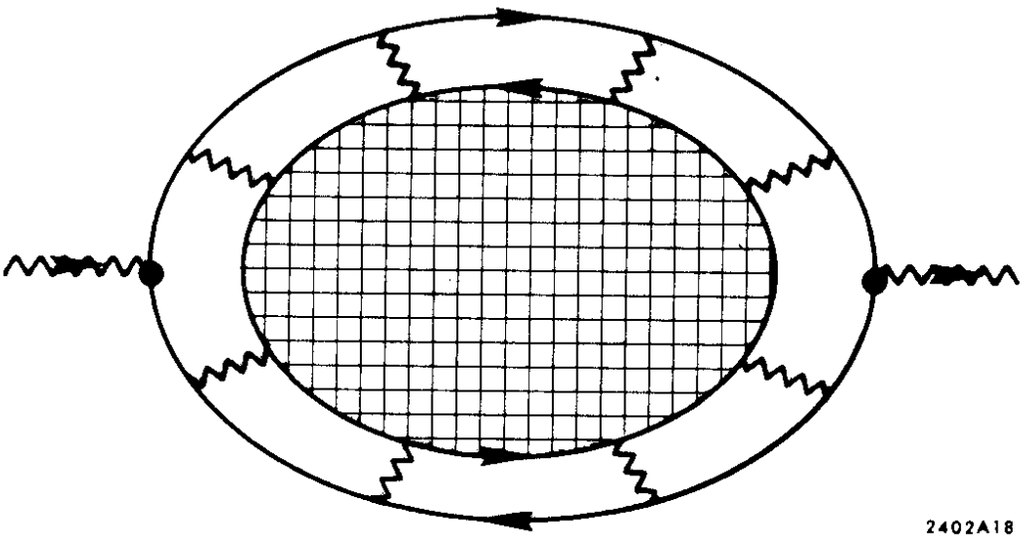

Fig. 8

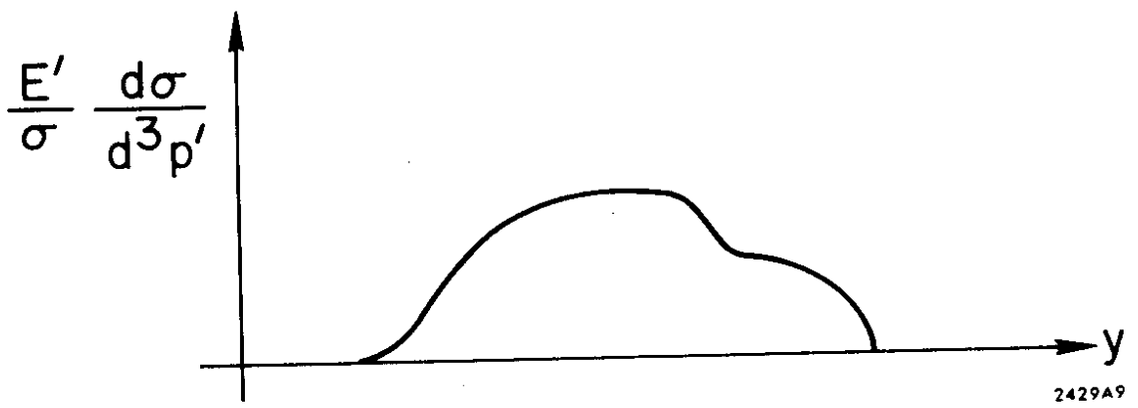

Fig. 9

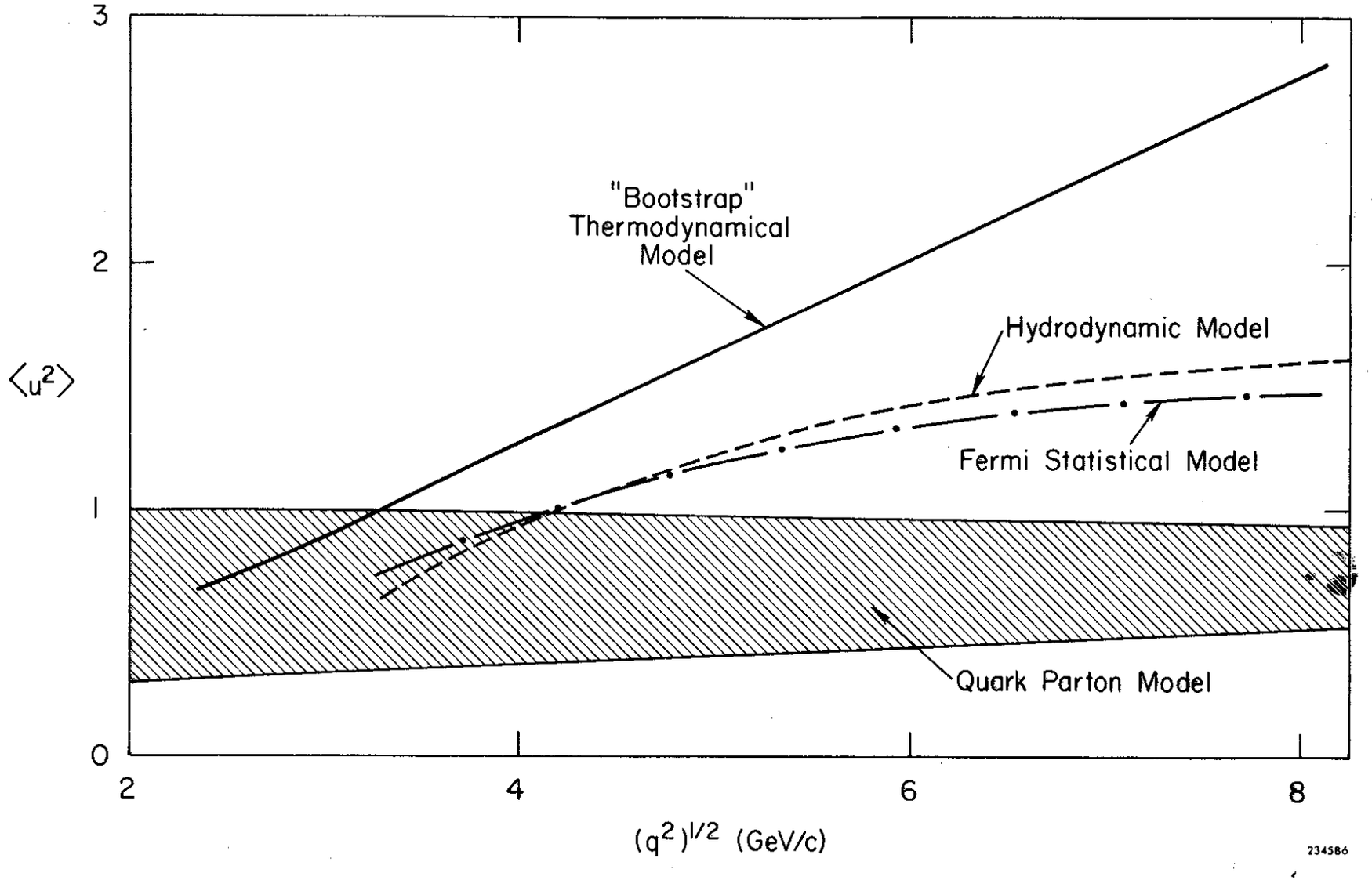
3
2
1
0
$\langle u^2 \rangle$
"Bootstrap"
Thermodynamical
Model
Hydrodynamic Model
Fermi Statistical Model
Quark Parton Model
2
4
6
8
$(q^2)^{1/2}$ (GeV/c)
234586

Fig. 10