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SUM RULES FOR THE STRUCTURE PARAMETERS OF THE NUCLEON

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A B S T R A C T

Based on the parton model and some commonly used assumptions, a class of sum rules is obtained which relates the structure parameters of the nucleon such as the proton charge radius with the quantities which can be most directly connected with the experimental observables in the hard massive photon limit, that is, the deep inelastic nucleon structure functions.

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The Cabibbo-Radicati sum rule ¹⁾ relates the isovector charge radius of the nucleon with the integrated form of the isovector on-shell photon-nucleon total cross-section. Its validity tests the basic assumption of the Gell-Mann local current algebra ²⁾ on which the derivation of the sum rule is essentially based. In this paper we shall obtain a class of sum rules which relate the structure parameters of the nucleon with the quantities which can be most directly connected with the experimental observables in the hard massive photon limit, that is, the deep inelastic structure functions of the nucleon. Among the sum rules there are not only those which relate the asymptotic parameters of the nucleon form factor for large q^2 to the threshold behaviour of the nucleon inelastic structure functions, but also one which expresses the nucleon charge radius at $q^2 = 0$ in terms of the integrated forms of the inelastic structure functions.

Our sum rules are based on the parton model proposed by Feynman ³⁾ and developed extensively by Bjorken and Paschos ⁴⁾ and Drell, Levy and Yan ⁵⁾. Within this framework we shall develop our study making the following specific but commonly used assumptions ^{*}) on the parton dynamics ; (i) the transverse cut-off of parton momenta and (ii) the symmetric distribution of partons inside the nucleon. We also assume (iii) the factorized form of the parton distribution function with respect to the longitudinal and the transverse momenta in the infinite momentum frame.

In order to transform these assumptions into an analytic form, we take the symmetric parton distribution function $f_P(x_1 \dots x_N ; \underline{k}_{1\perp} \dots \underline{k}_{N\perp})$ ^{**)} to be

$$f_P(x_1 \dots x_N ; \underline{k}_{1\perp} \dots \underline{k}_{N\perp}) = f_N(x_1 \dots x_N) g_N(\underline{k}_{1\perp} \dots \underline{k}_{N\perp}) \quad (1)$$

^{*}) We believe that some of the assumptions to be made in the following such as (ii) could be relaxed although it might cause some complications in expressing the final answers.

^{**)} $f_P(x_1 \dots x_N ; \underline{k}_{1\perp} \dots \underline{k}_{N\perp})$ is defined in the infinite momentum frame as the transition amplitude from the physical nucleon to the N parton state with the longitudinal and transverse parton momenta specified by $(x_1 \dots x_N ; \underline{k}_{1\perp} \dots \underline{k}_{N\perp})$. The suffix P specifies the longitudinal direction. For the detailed definition, see Drell and Yan, Ref. 6).

with the Gaussian cut-off ^{*}) for the transverse momenta $\underline{k}_{i\perp}$:

$$g_N(\underline{k}_{1\perp} \dots \underline{k}_{N\perp}) = \mathcal{N} \prod_{i=1}^N \sqrt{\frac{1}{\pi \langle \underline{k}_\perp^2 \rangle}} \exp\left(-\frac{\underline{k}_{i\perp}^2}{2 \langle \underline{k}_\perp^2 \rangle}\right) \quad (2)$$

In the above x_i represents the longitudinal fraction of the total momentum P of the system carried by the i th parton. We choose

$$P_\mu = (0, 0, P, P + \frac{M^2}{2P}) , \quad \underline{k}_{i\perp} \cdot \underline{P} = 0$$

with nucleon mass M and the infinite momentum frame is defined by the limit $P \rightarrow \infty$. We have the normalization condition

$$\sum_N \int \prod_{i=1}^N dx_i d^2 \underline{k}_{i\perp} \delta(1 - \sum_{i=1}^N x_i) \delta^2(\sum_{i=1}^N \underline{k}_{i\perp}) \cdot f_{P'}^*(x_1 \dots x_N; \underline{k}_{1\perp} \dots \underline{k}_{N\perp}) f_P(x_1 \dots x_N; \underline{k}_{1\perp} \dots \underline{k}_{N\perp}) = \delta^3(\underline{P}' - \underline{P}) \quad (3)$$

We shall impose the following normalization on g_N :

$$\int \prod_{i=1}^N d^2 \underline{k}_{i\perp} \delta^2(\sum_{i=1}^N \underline{k}_{i\perp}) |g_N(\underline{k}_{1\perp} \dots \underline{k}_{N\perp})|^2 = 1 \quad (4)$$

from which we obtain

$$|\mathcal{N}| = \sqrt{\pi N \langle \underline{k}_\perp^2 \rangle} \quad (5)$$

Now let us define $P(N)$ and $\mathcal{F}_N(x_1 \dots x_N)$ by

$$P(N) \mathcal{F}_N(x_1 \dots x_N) = |f_N(x_1 \dots x_N)|^2 \quad (6)$$

such that

$$\sum_N P(N) = 1 \quad (7)$$

$$\int_0^1 \prod_{i=1}^N dx_i \delta(1 - \sum_{i=1}^N x_i) \mathcal{F}_N(x_1 \dots x_N) = 1 \quad (8)$$

^{*}) For simplicity we assume this form. We could also work with a slightly generalized form of $\exp(-\sum_{i,j} C_{ij}^N \underline{k}_{i\perp} \cdot \underline{k}_{j\perp})$ where C_{ij}^N is a positive definite symmetric matrix. In general this form is a good representation for small $\underline{k}_{i\perp}$, since by making a Taylor expansion about $\underline{k}_{i\perp} = 0$ and keeping only the lowest order terms we see that $g_N(\underline{k}_{1\perp} \dots \underline{k}_{N\perp}) = \exp[-\ln(1/g_N)] \cong C_0 \exp(-\sum_{i,j} C_{ij}^N \underline{k}_{i\perp} \cdot \underline{k}_{j\perp})$ where the term linear in $\underline{k}_{i\perp}$ is absent due to rotational and parity invariance.

Then we find that $P(N)$ is the probability of finding a configuration of N partons in the nucleon and $\mathcal{F}_N(x_1 \dots x_N)$ is the joint probability of finding in such configurations partons with longitudinal fractions $(x_1 \dots x_N)$.

In the parton model, as Drell and Yan have shown ⁶⁾, the elastic electromagnetic form factor $F_1(q^2)$ of the nucleon is given in terms of f_P as

$$F_1(q^2) = \sum_{N=2}^{\infty} \int \prod_{i=1}^N dx_i d^2 \underline{k}_{i\perp} \delta(1 - \sum_{i=1}^N x_i) \delta^2(\sum_{i=1}^N \underline{k}_{i\perp}) \sum_a Q_a \cdot$$

$$f_P^*(x_1 \dots x_N; \underline{k}_{1\perp} - x_1 \underline{q}, \dots, \underline{k}_{a\perp} + (1-x_a) \underline{q}, \dots, \underline{k}_{N\perp} - x_N \underline{q}) f_P(x_1 \dots x_N; \underline{k}_{1\perp} \dots \underline{k}_{N\perp}) \quad , \quad (9)$$

where the infinite momentum frame is chosen such that

$$\underline{q}_\mu = (\underline{q}_\perp, 0, \frac{M\nu}{P}) \quad , \quad \nu = - \frac{P \cdot \underline{q}}{M}$$

and

$$q^2 = |\underline{q}_\perp|^2 + O(1/P^2) \quad .$$

In Eq. (9) the summation starts at $N = 2$, i.e., at the two-parton state ^{*)} in order to exclude the core contribution to the elastic form factor ^{**) .} Equation (9) can be understood easily in terms of the following intuitive picture (see the Figure). The nucleon with momentum P dissociates into multi-parton states with the transition amplitude f_P . Then a point-like parton a with charge Q_a interacts with the photon gaining the momentum q carried by the photon. The attained multi-parton configuration recombines itself back to the nucleon with added momentum $P' = P + q$. The amplitude $f_{P'}$ for the rejoining is known from f_P just by making a rotation from the direction $\underline{P}$ to $\underline{P}' = \underline{P} + \underline{q}$. This rotation alters longitudinal projection x_i only by corrective terms $O(1/P^2)$, whereas each momentum $\underline{k}_{i\perp}$ transverse to $\underline{P}$ is identical to order $1/P$ with $\underline{k}_{i\perp} - x_i \underline{q}$ as reckoned relative to $\underline{P}'$. Only the constituent a has its momentum altered by $\underline{q}$ as a result of the scattering by the photon current.

We point out that within the framework of the parton model the expression given by Eq. (9) is valid for any q^2 since there is only a single electromagnetic current operating there and all the strong interaction effects

^{*)} In the quark parton model the summation starts at $N = 3$.

^{**) This guarantees that the form factor vanishes as $q^2 \rightarrow \infty$. See Ref. 6).}

surviving in the infinite momentum limit is included through f_p and $f_{p'}$. On the other hand, the expression for the structure functions W_1 and νW_2 as is established in the papers of Refs. 4) and 5) and written as an incoherent sum of each parton contribution ($\omega = 2M\nu/q^2$)

$$(\nu W_2)_a^N = Q_a^2 \frac{1}{\omega} \int \prod_{i=1}^N dx_i d^2k_{i\perp} \delta(x_a - \frac{1}{\omega}) \delta(1 - \sum_{i=1}^N x_i) \delta^2(\sum_{i=1}^N k_{i\perp}) \left| f_p(x_1 \dots x_N; k_{1\perp} \dots k_{N\perp}) \right|^2 \quad (10)$$

is valid only in the Bjorken limit, $\nu, q^2 \rightarrow \infty$ with ω fixed. This is because only in this limit the free multi-parton states can saturate between the two "undressed" currents operating in the matrix elements of W_1 and νW_2 *).

Substituting Eq. (1) into Eq. (10) and using Eqs. (4) and (6), we obtain the familiar parton model expression for νW_2

$$\nu W_2(\omega) = \sum_{N=2}^{\infty} P(N) \langle \sum_a Q_a^2 \rangle_N \frac{1}{\omega} \mathcal{F}_N(\frac{1}{\omega}) \quad (11)$$

where

$$\mathcal{F}_N(x_i) = \int \prod_{i=2}^N dx_i \delta(1 - \sum_{i=1}^N x_i) \mathcal{F}_N(x_1 \dots x_N) \quad (12)$$

and $\langle \sum_a Q_a^2 \rangle_N$ equals the average value of $\sum_a Q_a^2$ in a configuration of N partons in the nucleon. We notice that the specific form of the transverse momentum cut-off does not play any crucial rôle in obtaining Eq. (11). Similarly we evaluate the elastic form factor by substituting Eqs. (1) and (2) into Eq. (9). Using the formula $\delta^2(\underline{k}) = (2\pi)^{-2} \int d^2b \exp(i\underline{k}\underline{b})$, one can make the $\underline{k}_{i\perp}$ integration easily. Then, performing the $\underline{b}$ integration, we finally get **)

*) More precisely, only in the Bjorken limit one can make the following substitution in the absorptive amplitude $W_{\mu\nu} : \langle P | J_\mu(x) J_\nu(0) | P \rangle \rightarrow \langle U(0)P | j_\mu(x) j_\nu(0) | U(0)P \rangle$ where $J_\mu(x)$ is the fully interacting electromagnetic current and $j_\mu(x)$ is the corresponding free current and both are related through the U matrix as $J_\mu(x) = U^{-1}(t) j_\mu(x) U(t)$. Refer to the first paper in Ref. 5).

**) A similar formula was also obtained by K. Kitani (Tokyo Institute of Technology Preprint, 1972). In his derivation, however, an incorrect approximation was made by unnecessarily relaxing the conservation of the transverse momenta. An importance of our exact treatment will become clear in the discussion which follows Eqs. (14) and (15).

$$F_1(q^2) = \sum_{N=2}^{\infty} P(N) \langle \sum_a Q_a \rangle_N \int \prod_{i=1}^N dx_i \delta(1 - \sum_{i=1}^N x_i) \mathcal{F}_N(x_1 \dots x_N) \exp \left[-\frac{q^2}{4 \langle k_{\perp}^2 \rangle} \left\{ (1-x_1)^2 + \sum_{i=2}^N x_i^2 \right\} \right] \quad (13)$$

In the derivation of Eqs. (11) and (13) we have used the assumption (ii) which allows to factor out the charge distributions $\langle \sum_a Q_a^2 \rangle_N$ and $\langle \sum_a Q_a \rangle_N$. From charge conservation $\langle \sum_a Q_a \rangle_N = 1$ for the proton and $\langle \sum_a Q_a \rangle_N = 0$ for the neutron. This predicts for the elastic form factor of the neutron that $F_1(q^2)_{\text{neutron}} \equiv 0$, which is consistent with the experiments ^{7) *}. In the following we only consider the proton form factor.

An interesting sum rule is obtained by taking the derivative with respect to q^2 in Eq. (13). Noting the relation $-F_1'(0)_{\text{proton}} = \langle r_1^2 \rangle_p / 6$, we obtain

$$\frac{2}{3} \langle r_1^2 \rangle_p = \frac{1}{\langle k_{\perp}^2 \rangle} \left[1 - \sum_{N=2}^{\infty} P(N) (2 \langle x \rangle_N - N \langle x^2 \rangle_N) \right] \quad (14)$$

where

$$\begin{aligned} \langle x \rangle_N &= \int_0^1 dx \, x \mathcal{F}_N(x) = \frac{1}{N} \quad , \\ \langle x^2 \rangle_N &= \int_0^1 dx \, x^2 \mathcal{F}_N(x) \end{aligned} \quad (15)$$

and Eq. (7) was used. This sum rule expresses the charge radius of the proton in terms of the moments of x and the transverse cut-off parameter $\langle k_{\perp}^2 \rangle$ of the parton momenta. Naturally one expects that there is a strong correlation between $\langle k_{\perp}^2 \rangle$ and the proton size $\langle r_1^2 \rangle_p$. The naïve application of the uncertainty relation tells us that $(2/3) \langle r_1^2 \rangle_p \approx 1 / \langle k_{\perp}^2 \rangle$, which actually corresponds to the contribution of the first term in the right-hand side of Eq. (17). What is the rest ?

Here a remark is in order. Consider a system of N independent particles with a Gaussian wave function of momentum fluctuation $(\Delta p_x)^2$ and impose total momentum conservation on the system. Then, the spacial spread

*) If a non-zero form factor for the neutron were experimentally favoured, it could be obtained by relaxing the assumption (ii).

given by $(\Delta x)^2 = 1/(\Delta p_x)^2$ for the single Gaussian wave function is changed due to the total momentum conservation of the many-particle system into the form $(\Delta x)^2 = (1 - 1/N)/(\Delta p_x)^2$. If each particle carries a charge Q_a , then the charge radius will be $(2/3) < r_1^2 >_N = (\sum_{a=1}^N Q_a)(1 - 1/N)/(\Delta p_x)^2$ where $\sum_{a=1}^N Q_a$ is the total charge of the system. In the parton model the proton can be regarded as composed of multi-partons interacting with each other in a space of the finite size given by $(\Delta x)^2 = (2/3) < r_1^2 >_p$. Then noting that $(\Delta p_x)^2 = < k_{\perp}^2 >$ we can rewrite the sum rule (14) as

$$\frac{2}{3} < r_1^2 >_p = \frac{1}{< k_{\perp}^2 >} \sum_{N=2}^{\infty} P(N) \left[\left(1 - \frac{1}{N}\right) - \left(\frac{1}{N} - N < x^2 >_N\right) \right], \quad (16)$$

where the relation (15) was used. The interpretation is obvious. The first term on the right-hand side is the contribution from the free parton fluctuation with the constraint of total momentum conservation. The second term represents the correction to the first contribution due to the interaction between the multi-partons. Dynamics enters also through the probability function $P(N)$. Apparently we expect that if the component of many-parton states is dominant in the composition of the proton, the correction of the second term may be small, whereas if the component of a few parton states is overwhelming, then the correction must be substantial and the detailed parton dynamics will be crucial for its estimation. Rather than attacking the problem directly, we shall here attempt to express the right-hand side of Eq. (14) in terms of the experimental observables.

For this purpose we introduce the specific assumption of (iv) the quark parton model which tells us the charge distribution of partons inside the nucleon. Assuming that in addition to the three quarks there is a statistical distribution of the quark-antiquark pairs inside the nucleon ^{*)}, we get ⁴⁾

$$\begin{aligned} \left\langle \sum_a Q_a^2 \right\rangle_N &= 1 + \frac{2}{9}(N-3) = \frac{2}{9}N + \frac{1}{3} \quad \text{for the proton,} \\ &= \frac{2}{3} + \frac{2}{9}(N-3) = \frac{2}{9}N \quad \text{for the neutron.} \end{aligned} \quad (17)$$

where $[(2/3)^2 + (1/3)^2 + (1/3)^2]/3 = 2/9$. Using these charge distributions in Eq. (11), we obtain

^{*)} There is an experimental indication ⁸⁾ that parton models require neutral gluons in order to reproduce the experimental numbers in their sum rules. The introduction of the gluons does not change Eq. (18) but modifies Eq. (19). However, we have found that because the available data, particularly for the neutron structure functions, are not sufficiently accurate, this modification due to the gluons is not so significant to our qualitative analysis in the following. I thank Dr. M. Gronau for discussions on this point.

$$\sum_N P(N) \langle x \rangle_N = 3 \int_1^{\omega_0} \frac{d\omega}{\omega^2} (\nu W_2^p - \nu W_2^n) \equiv 3 I_1 \quad (18)$$

$$\sum_N P(N) N \langle x^2 \rangle_N = \frac{9}{2} \int_1^{\omega_0} \frac{d\omega}{\omega^3} \nu W_2^x \equiv \frac{9}{2} I_2 \quad (19)$$

Therefore, from Eqs. (14), (18) and (19) we finally obtain

$$\frac{2}{3} \langle r_1^2 \rangle_p = \frac{1}{\langle k_\perp^2 \rangle} \left(1 - 6 I_1 + \frac{9}{2} I_2 \right) \quad (20)$$

We repeat that the second and third terms on the right-hand side are the dynamical corrections of multi-parton states to the naïve consideration about the transverse momentum cut-off and the proton size. The importance is that they are expressed in terms of the measurable quantities.

The experimental numbers from the SLAC-MIT experiment ⁸⁾ are $3 I_1 \approx 0.172$, $I_2 \approx 0.025$. Therefore we have

$$\frac{2}{3} \langle r_1^2 \rangle_p = \frac{1}{\langle k_\perp^2 \rangle} (1 - 0.34 + 0.11) = \frac{0.77}{\langle k_\perp^2 \rangle} \quad (21)$$

In view of the discussion which follows Eq. (16) we rewrite the factor in the right-hand side of Eq. (20) as $[1 - 3 I_1 - (3 I_1 - 9 I_2/2)]$ and put the numbers for I_1 and I_2 . Then we see roughly that out of the total correction 23%, 17% is explained in terms of the correction due to the momentum conservation constraint on the independent particle fluctuation, whereas the remaining 6% is due to the detailed dynamics of partons. We must admit that the detailed numbers are heavily dependent on the quark parton model, which we adopted in order to make the explicit numerical estimation. However, the following conclusion which we draw from the consideration following Eq. (16) and the above numerical estimate must be stable : the component of many-parton states is dominant in the composition of the proton. Using the experimental number $\langle r_1^2 \rangle_p / 6 = 2.31 \text{ GeV}^{-2}$ obtained from the dipole fit to $G_E^p(q^2)$ ⁷⁾ in the above equality, we obtain $\sqrt{\langle k_\perp^2 \rangle} \approx 300 \text{ MeV}$. This seems to be reasonable.

It is a straightforward application to obtain other sum rules by taking higher derivatives of $F_1(q^2)_{\text{proton}}$ at $q^2 = 0$ and expressing them in terms of the corresponding higher moments of x . But the physical interpretation seems to be less straightforward, so we shall not discuss it here.

Now, let us discuss the relations for the asymptotic parameters of the proton form factor. Define a function

$$G(q^2) = \int_0^1 dx \sum_{N=2}^{\infty} P(N) \mathcal{F}_N(x) \exp\left[-\frac{q^2}{4\langle k_{\perp}^2 \rangle} (1-x_1)^2\right] \quad (22)$$

Then using the inequalities $(1-x_1)^2 \leq (1-x_1)^2 + (1-x_1)^2/(N-1)$, $(1-x_1)^2 + \sum_{i=2}^N x_i^2 \leq 2(1-x_1)^2$, it is straightforward to prove

$$G(2q^2) \leq F_1(q^2)_{\text{proton}} \leq G(q^2) \quad (23)$$

Again based on the quark parton model we can express $G(q^2)$ in terms of the measurable quantity through the relation $\sum_N P(N) \mathcal{F}_N(x) = 3[\nu W_2^p(1/x) - \nu W_2^n(1/x)]$. Suppose that the asymptotic behaviours are given by

$$[\nu W_2^p(\omega) - \nu W_2^n(\omega)] \sim A^{\frac{p+1}{2}} \left(1 - \frac{1}{\omega}\right)^p \text{ as } \omega \rightarrow 1, \quad (24)$$

$$F_1(q^2)_{\text{proton}} \sim B \left(\frac{0.71}{q^2}\right)^{-n/2} \text{ as } q^2 \rightarrow \infty, \quad (25)$$

then from Eq. (22) we are able to estimate

$$G(q^2) \sim \frac{3}{2} T\left(\frac{p+1}{2}\right) \left(\frac{4A\langle k_{\perp}^2 \rangle}{q^2}\right)^{\frac{p+1}{2}} \text{ as } q^2 \rightarrow \infty. \quad (26)$$

From Eqs. (23), (25) and (26) we immediately confirm the Drell-Yan relation ⁶⁾

$$n = p + 1 \quad (27)$$

We can also obtain the inequality relation between A , B and $\langle k_{\perp}^2 \rangle$. Here, in order to save space and to make the numerical check more concrete, we take $n = 4$ from the dipole fit of the form factors ⁷⁾. Then $p = 3$ from the above relation (27) and we obtain

$$\frac{0.71 \sqrt{B}}{2\sqrt{6} A} \leq \langle k_{\perp}^2 \rangle \leq \frac{0.71 \sqrt{B}}{\sqrt{6} A} \quad (28)$$

Substituting the experimentally suggested numbers $A^2 = 1.3$ from the SLAC-MIT experiment ^{9),8)} and $B = 1 + \kappa_p = 2.79$ from the dipole fit ⁷⁾, we obtain

$$460 \leq \sqrt{\langle k_{\perp}^2 \rangle} \leq 650 \text{ in MeV/c} \quad (29)$$

The answer seems to be of right order. We remark that the resulting numbers are not to be taken that seriously since for example A^2 is experimentally the poorly determined quantity (the value of A^2 might be as small as 0.65 ⁹⁾).

Finally we mention that the cut-off parameter $\langle k_{\perp}^2 \rangle$ for the parton transverse momentum has no direct comparison with the experimentally observed transverse momentum cut-off ^{*)} of produced particles in high energy reactions ¹⁰⁾. However, we expect that the two are closely related to each other and further that in the solution of hadron dynamics they will come out to be of the same order of magnitude. Our results confirm this guess.

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^{*)} Experimentally we observe $\langle p_{\perp} \rangle = 300 \sim 400 \text{ MeV/c}$ ¹⁰⁾. In general we have $\sqrt{\langle p_{\perp}^2 \rangle} \gtrsim \langle p_{\perp} \rangle$.

REFERENCES

- 1) N. Cabibbo and L.A. Radicati - Phys.Letters 19, 697 (1966).
- 2) S.L. Adler and R.F. Dashen - "Current Algebra and Applications to Particle Physics", W.A. Benjamin, Inc., New York (1968).
- 3) R.P. Feynman - Phys.Rev.Letters 23, 1415 (1969) ; Proceedings of the Third Topical Conference on High Energy Collisions of Hadrons, Stony Brook, New York (1969), p. 237.
- 4) J.D. Bjorken and E.A. Paschos - Phys.Rev. 185, 1975 (1969).
- 5) S.D. Drell, D.J. Levy and T.M. Yan - Phys.Rev. 187, 2159 (1969) ; *ibid* D1, 1617 (1970) ; *ibid* D1, 1035 (1970) ;
T.M. Yan and S.D. Drell - Phys.Rev. D1, 2402 (1970) ;
S.D. Drell and T.M. Yan - Ann.Phys.(N.Y.) 66, 578 (1971).
- 6) S.D. Drell and T.M. Yan - Phys.Rev.Letters 24, 181 (1970).
- 7) R. Wilson - Proceedings of the 1971 International Symposium on Electron and Photon Interactions at High Energies, Ithaca, New York (1971), p. 97.
- 8) H.W. Kendall - Proceedings of the 1971 International Symposium on Electron and Photon Interactions at High Energies, Ithaca, New York (1971), p. 247. Also, MIT-SLAC Report, XV International Conference on High Energy Physics, Kiev, U.S.S.R. (1970).
- 9) SLAC-MIT Report, SLAC-PUB-907 (1971).
- 10) See, for example :
D.R.O. Morrison - Review Talk given at the Fourth International Conference on High Energy Collisions, Oxford (April 1972).

FIGURE CAPTION

Parton diagram for the elastic form factor of the nucleon.

FIG.1