Here $\Sigma(x)$ comes from the classical vortex solution in the $A_0=0$ gauge.\(^4\)

Finally, I wish to remark that we have done a matching between two different perturbations schemes; namely, between the standard perturbation and the one around the instanton. This is a new concept in field theory which may have other applications besides tunneling as Gribov's problem.\(^5\)

**References**


6. C. N. Yang: talk given in Session C2, this Conference.

7. C. Callan: talk given in Session C1, this Conference.
b) The central technical point is that there exists a well defined procedure to realize the program. Namely, one can rely on the standard operator expansion (O.E.), to treat the first power corrections in $Q^2$.

\[ \int e^{i s \tau} T \{ j_\mu (x) j_\nu (0) \} d^4 x = (q_\mu q_\nu - g_\nu q^2) \sum_n O_n \]

the few lower dimension operators being $O_n = I, G_\mu^a G_\nu^a, m_\mu \bar{\psi} \gamma_\mu \Gamma \phi \bar{\psi} \Gamma \phi$ where $G_\mu^a$ is the gluon field strength, $\phi$ is a quark field and $\Gamma$ are some matrices.

O.E. is so widely used nowadays that it might look strange to dwell upon it here. But in fact all the discussion so far refers to the perturbation theory where one has:

1) O.E. is OK.
2) Taking vacuum average of $\sum_n O_n$ picks up the unity operator.

We are interested in nonperturbative effects of QCD, however, and both statements are modified:

1') O.E. is meaningful only as far as the corresponding matrix elements are dominated by the large scale, i.e., $Q^2$ independent fluctuations.

2') Apart from the unity operator the others develop nonvanishing vacuum expectation as well ($\langle 0 | G_\mu^a G_\nu^a | 0 \rangle, \langle 0 | \bar{\psi} \Gamma \phi \bar{\psi} \Gamma \phi | 0 \rangle \cdots \neq 0$).

Fortunately enough, we do know that at short distances the leading nonperturbative effects are given by instantons and are able to justify the use of operator expansion for up to quite a high power of $(Q^{-2})$.

c) In fact we use a new form of sum rules which comes out by applying some operator $\hat{L}$ both to the dispersion and QCD representations for the polarization operators. The result looks as

\[ \int e^{-s/M^2} \sigma_{\text{phys}}(s) - \sigma_{\text{quark}}(s) ds = h_2 \frac{M^4}{M^6} + \frac{1}{M^6} + O(M^{-8}), \]

where the coefficients $h_{2,5}$ are given in terms of $\langle 0 | G^a_\mu G^a_\nu | 0 \rangle, \langle 0 | \phi^4 | 0 \rangle$. Note the appearance of a sharp exponential weight in the integrals over the cross sections. The precise form of operator $\hat{L}$ is

\[ \hat{L} = \lim_{n \to \infty} \frac{1}{n!} \left( Q^2 dQ^2 \right)^n, \]

with $Q^2 = M^2, Q^2, n \to \infty$.

It is amusing to find that applying the operator $\hat{L}$ is equivalent to the Borel summation of the power corrections. Thus we suppress higher powers in $Q^{-2}$ which are neglected. The procedure can be reiterated but the second Borel transform would introduce a sign alternating weight function. Therefore the first Borel transform is the best one to study resonances.

d) Applications of the sum rules are quite numerous since we can try different currents $J_\mu$. Our claim is that power corrections become essential only at a rather low mass $M^2 \approx m_\rho^2$ so that integral over $\sigma_{\text{phys}}(s)$ is saturated by a single resonance. Once we get power corrections of order 0.1-0.2 of the integral over the quark cross section we stop the calculation as not to go into the mass of higher order terms in $M^{-2}$.

In fact, for a well defined reason terms $\sim M^{-4}, M^{-6}$ become comparable at this point while higher powers of $M^{-2}$ are still hopefully much smaller.

Thus, quantitative results require two parameters $\langle 0 | G_\mu^a G_\nu^a | 0 \rangle, \langle 0 | \bar{\psi} \Gamma \phi \bar{\psi} \Gamma \phi | 0 \rangle$. The latter one can be in fact borrowed from the studies of the chiral symmetry breaking so that there is no problem. As for the former one we see no way but to sacrifice one of the sum rules (for charmonium decays) to fix it from the data:

\[ \langle 0 \mid \frac{\alpha_s}{\pi} G_\mu^a G_\nu^a \mid 0 \rangle \simeq 0.012 \text{ GeV}^4. \]

Then we are free to evaluate the masses, mixing angles and leptonic widths of $\rho, \omega, \phi$. The axial vector channel (the $\pi - A_1$ system) is treated in the same way. Changing the current modifies only the coefficients $h_{2,5}$ in a well defined way and it is a real amusement to see that the resonance properties do comply with the change in the sum rules.
The analysis seems to have some implications for the other models. In particular, via the triangle anomaly in the trace of the energy-momentum tensor $\langle 0 | \alpha_s G^2 | 0 \rangle$ is related to the vacuum energy due to the nonperturbative effects. This is the energy $\varepsilon$ outside the bag. Quarks most probably destroy vacuum fluctuations and, therefore, we propose an estimate for the constant $B$ of the bag model:

$$B = -\varepsilon = -\frac{9}{32} \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu
u} G^{a\mu\nu} | 0 \rangle$$

$$= -0.0035 \text{ GeV}^2$$

which is not bad numerically. Further argument can be found in the paper by Shuryak.\(^3\)

References
Session C5: Formal Theory

Chairman: A. Martin
Organizer: K. Nishijima
Scientific Secretaries: M. Kobayashi
H. Nakajima

1. Manifestly Covariant Canonical Formalism of Non-Abelian Gauge Theory and Quantum Gravity
   N. Nakanishi

2. Relativistic Nonlinear Quantum Mechanics
   T. W. B. Kibble

3. Asymptotic Freedom and the Symplectic and $G_2$ Groups
   M. Chachian

(Friday, August 25, 1978; 9:00–10:20)