Title
LOCAL POUR QUARK OPERATOR IN K K MIXING: THE VACUUM SATURATION ESTIMATE AS AN UPPER BOUND FOR THE MATRIX ELEMENT

Permalink
https://escholarship.org/uc/item/62r4t49g

Author
Machet, B.

Publication Date
1984-04-01

LBL-17728
Preprint c.2

Lawrence Berkeley Laboratory

UNIVERSITY OF CALIFORNIA

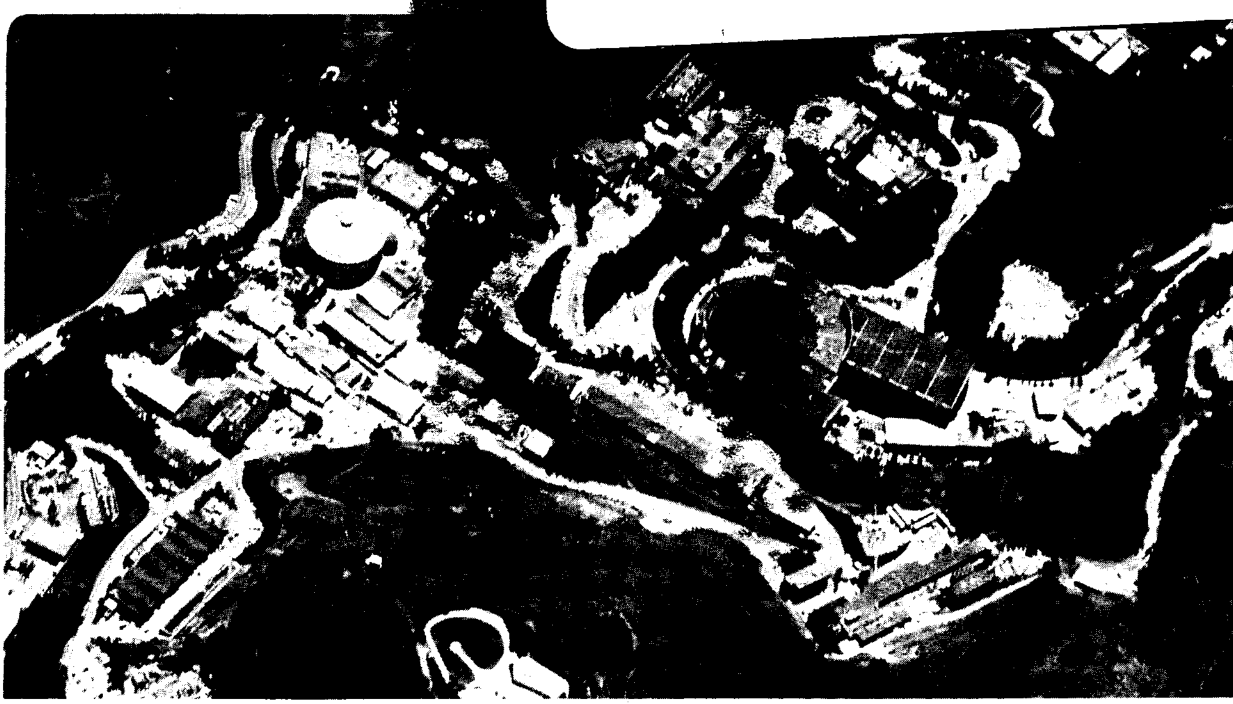
Physics Division

Submitted for publication

LOCAL FOUR QUARK OPERATOR IN $K^0 K^0$ MIXING: THE VACUUM SATURATION ESTIMATE AS AN UPPER BOUND FOR THE MATRIX ELEMENT

B. Machet

April 1984



Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

DISCLAIMER

April 1984 LBL-17728

LOCAL FOUR QUARK OPERATOR IN K^0K^0 MIXING: THE VACUUM SATURATION ESTIMATE AS AN UPPER BOUND FOR THE MATRIX ELEMENT[†]

B. Machet*

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

Abstract

By computing the imaginary part of the propagator of the local four quark operator $O(x) = s_L(x)\gamma^\mu d_L(x) s_L(x)\gamma_\mu d_L(x)$, we obtain the upper bound:

$$| \quad | \equiv \langle K^0|O|K^0\rangle \leq | \quad |_{\text{vacuum estimate}} + O(1/N_c^2)$$

with an uncertainty (outside the $1/N_c^2$ corrections) of order 10%. We show how it can be recovered by Laplace hadronic sum rules.

[†] *This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.*

**Participating guest at Lawrence Berkeley Laboratory*
On leave from: CNRS, Centre de Physique Théorique, Luminy, Case 907, F13288 Marseille Cédex 9, FRANCE

I. INTRODUCTION

Weak interactions generate a mixing between the K^0 and $\bar{K}^0$ states,[1] mainly by the exchange of two W bosons (Fig. 1). (We neglect the Higgs contributions.)[2] In a free quark model, this transition can be considered to be mediated by an effective four quark local hamiltonian;[1] this property subsists when one resums all hard gluonic corrections at the leading logarithm approximation.[2,3] We shall deduce an upper bound for the matrix element:

$$|\mathcal{M}| = |\langle \bar{K}^0 | O | K^0 \rangle| \tag{1}$$

of this <u>local</u> operator

$$O(x) = \bar{s}_L^a(x)\gamma^\mu d_L^a(x)\,\bar{s}_L^b(x)\gamma_\mu d_L^b(x) \;. \tag{2}$$

(a and b are color indices).

It is based on the computation of the imaginary part of the correlation function:

$$\psi(q^2) = i\int d^4x\, e^{iq\cdot x}\langle 0|T\; O(x)\; O^+(0)|0\rangle \tag{3}$$

up to corrections of order $1/N_c^2$.

The result, very close to the vacuum saturation estimate[1,3]

$$|\mathcal{M}|_{vac} = \frac{4}{3} f_K^2 M_K^2 \;, \tag{4}$$

strengthens the constraint obtained recently[4] by hadronic sum rules, can be simply recovered in their "Laplace" (or "Borel") version[5] and constitutes in this precise case their maximal capability.

II. DEDUCTION OF THE BOUND. DISCUSSION

We shall be concerned henceforward with the matrix element of the local four quark operator O(x). This means in particular that we deliberately ignore the contributions of soft gluons (Fig. 2a), light quark condensates (Fig. 2b) etc. ...,which we cannot see how they can be handled within this approximation.

If one neglects the mixing angles other than Cabibbo's θ_c, the K_L-K_S mass splitting writes:

$$M_L - M_S = \frac{G_F^2}{4\pi^2} \sin^2\theta_c \cos^2\theta_c \, m_c^2 \, \eta \, \frac{\mathcal{M}}{M_{K^0}} \quad , \tag{5}$$

with $\mathcal{M}$ defined in Eqs. (1,2); η originates from the resummation of hard gluons corrections at the leading logarithm approximation.[2,3]

Looking for an upper bound for the modulus squared $|\mathcal{M}|^2$, we study the imaginary part of the corrrelation function $\psi(q^2)$ defined in Eq. (3), graphically depicted in Fig. (3). At order $1/N_c^2$ [6] it factorizes into Fig. (4): gluonic corrections to Fig. (4) breaking the factorization are indeed at least suppressed by 2 powers of $1/N_c$. For example, Fig. (5b) = $O(1/N_c^2)$ Fig. (5a) (2 gluons at least are needed by Furry's theorem), Fig. (5d) = $O(1/N_c^2)$ Fig. (5c), etc... . The analytic expression corresponding to Fig. (4) is readily obtained by separately resumming at all orders of the strong interactions the hadronic current propagators appearing at the free quark level in the computation of Fig. (6a), leading to:

$$\psi(q^2) = -2i \, \frac{1}{16} \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4 \ell}{(2\pi)^4} (2\pi)^4 \delta^4(q-k-\ell)$$

.../...

$$\left[\Pi^{\mu\nu\, 6+i7}_{VV}(k) + \Pi^{\mu\nu\, 6+i7}_{AA}(k) \right] \left[\Pi^{6+i7}_{VV\, \mu\nu}(\ell) + \Pi^{6+i7}_{AA\, \mu\nu}(\ell) \right] . \tag{6}$$

The $\Pi_{\mu\nu}$'s are the propagators of the vector (V V) or axial (AA) hadronic currrents

$$V_\mu^{6+i7}(x) = \bar d(x) \gamma_\mu s(x) \quad ,$$
$$A_\mu^{6+i7}(x) = \bar d(x) \gamma_\mu \gamma_5 s(x) \quad , \tag{7}$$

defined as:

$$\Pi^{\mu\nu\, 6+i7}_{VV\,(AA)}(q) = i \int d^4x \, e^{iq\cdot x} \langle 0 | T \; V^\mu(x) \, V^{\nu +}(0) \; (A^\mu(x) \, A^{\nu +}(0)) | 0 \rangle \; . \tag{8}$$

The same trick can be applied (up to corrections in $1/N_c^2$) to a formal resummation of Fig. (6b) where the fermionic lines of color have been crossed with respect to Fig. (6a). Fig. (6b) has the same analytic expression in terms of the $\Pi^{\mu\nu}$'s as Fig. (6a), up to an extra factor $1/N_c$. This leads to a factor $(1 + 1/N_c) = 4/3$ in Eq. (6) and to the final expression valid up to corrections of order $1/N_c^2$:

$$\psi(q^2) = -2i \, \frac{1}{16} \left(1 + \frac{1}{3}\right) \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4\ell}{(2\pi)^4} \, \delta^4(q-k-\ell)$$
$$\left[\Pi^{\mu\nu\, 6+i7}_{VV}(k) + \Pi^{\mu\nu\, 6+i7}_{AA}(k) \right] \left[\Pi^{6+i7}_{VV\, \mu\nu}(\ell) + \Pi^{6+i7}_{AA\, \mu\nu}(\ell) \right] . \tag{9}$$

The matrix element $\mathcal{M}$ can be considered[4] as the value at zero transfer (t = 0) of a scalar form factor F(t) analytic real in the complex t plane with a cut from $t_{0KK} = 4M_K^2$ to ∞. Its modulus squared $|F(t)|^2$ is directly related to Im $\psi(t)$, $|K^0 K^0>$ being the lowest intermediate state contributing to the absorptive part of ψ. This leads[4] to the inequality:

$$\frac{1}{\pi}\,\mathrm{Im}\,\psi(t) \geqslant \frac{1}{16\pi^2}\,\frac{1}{2}\sqrt{\frac{t-t_{0_{KK}}}{t}}\,|F(t)|^2\,\theta(t-t_{0_{KK}})\,.$$

$$t_{0_{KK}} = (2M_K)^2 \,. \tag{10}$$

From Eq. (9), on the other side, Im $\psi(t)$ can be estimated as a sum of contributions from low energy resonances and poles, plus a continuum which can be obtained from the asymptotic behavior of QCD (free quarks). In particular, as will be shown below, we can obtain a very simple analytic expression in the narrow width approximation, the relevance of which will be discussed a little further.

Limiting ourselves to the lowest (K, K*, Q) contributions and using CVC, we have the relations:

$$\frac{1}{\pi}\,\mathrm{Im}\,\overset{\text{6ait}}{\Pi}{}^{(0)}_{VV}(t) = 0 \qquad \text{(a)}$$

$$\frac{1}{\pi}\,\mathrm{Im}\,\overset{\text{6ait}}{\Pi}{}^{(1)}_{VV}(t) = \frac{M_{K^*}^2}{g_{K^*}^2}\,\delta(t-M_{K^*}^2) \qquad \text{(b)}$$

$$\frac{1}{\pi}\,\mathrm{Im}\,\overset{\text{6ait}}{\Pi}{}^{(0)}_{AA}(t) = 2f_K^2\,\delta(t-M_K^2) \qquad \text{(c)} \tag{11}$$

$$\frac{1}{\pi}\,\mathrm{Im}\,\overset{\text{6ait}}{\Pi}{}^{(1)}_{AA}(t) = \frac{M_Q^2}{g_Q^2}\,\delta(t-M_Q^2) \qquad \text{(d)}$$

for the functions $\Pi^{(0)}$ and $\Pi^{(1)}$ appearing in the decomposition of the $\Pi^{\mu\nu}$ orthogonal in the J = 0 and J = 1 channels:

$$\Pi^{\mu\nu}(q) = -(g^{\mu\nu}q^2 - q^\mu q^\nu)\,\Pi^{(1)}(q^2) + q^\mu q^\nu\,\Pi^{(0)}(q^2)\,. \tag{12}$$

Evaluating dispersively the $\Pi^{\mu\nu}$'s in Eq. (9) and integrating immediately over the dispersive variables t and t' we get:

$$\psi(q^2) = -2i\,\frac{1}{16}\left(1+\frac{1}{3}\right)\int\frac{d^4k}{(2\pi)^4}\int\frac{d^4p}{(2\pi)^4}\,(2\pi)^4\,\delta^4(q-k-p)$$

$$\Big\{\left[2k^2p^2+(k.p)^2\right]\Big[\left(\frac{M_{K^*}^2}{g_{K^*}^2}\right)^2\frac{1}{k^2-M_{K^*}^2}\,\frac{1}{p^2-M_{K^*}^2}+\left(\frac{M_Q^2}{g_Q^2}\right)^2\frac{1}{k^2-M_Q^2}\,\frac{1}{p^2-M_Q^2}$$

$$+\frac{M_{K^*}^2}{g_{K^*}^2}\,\frac{M_Q^2}{g_Q^2}\left(\frac{1}{k^2-M_{K^*}^2}\,\frac{1}{p^2-M_Q^2}+\frac{1}{k^2-M_Q^2}\,\frac{1}{p^2-M_{K^*}^2}\right)\Big]$$

$$+(k.p)^2\,(2f_K^2)^2\,\frac{1}{k^2-M_K^2}\,\frac{1}{p^2-M_K^2} \tag{13}$$

$$-\left[k^2p^2-(k.p)^2\right]2f_K^2\Big[\frac{M_{K^*}^2}{g_{K^*}^2}\left(\frac{1}{k^2-M_{K^*}^2}\,\frac{1}{p^2-M_K^2}+\frac{1}{k^2-M_K^2}\,\frac{1}{p^2-M_{K^*}^2}\right)$$

$$+\frac{M_Q^2}{g_Q^2}\left(\frac{1}{k^2-M_Q^2}\,\frac{1}{p^2-M_K^2}+\frac{1}{k^2-M_K^2}\,\frac{1}{p^2-M_Q^2}\right)\Big]\Big\}\,.$$

2 Im$\psi(t)$ is now obtained by replacing every term $1/(p^2-M^2)$ by the corresponding $2i\pi\,\delta(p^2-M^2)$ and performing explicatively the 2-body phase space integrals:

$$\mathrm{Im}\psi(t) = \frac{1}{16}\,\frac{1}{128\pi}\,\frac{1}{t}$$

$$\Big\{\left(\frac{M_{K^*}^2}{g_{K^*}^2}\right)^2\sqrt{t(t-t_{0_{K^*K^*}})}\left[3t_{0_{K^*K^*}}^2+4t(t-t_{0_{K^*K^*}})\right]\theta(t-t_{0_{K^*K^*}})$$

$$+\left(\frac{M_Q^2}{g_Q^2}\right)^2\sqrt{t(t-t_{0_{QQ}})}\left[3t_{0_{QQ}}^2+4t(t-t_{0_{QQ}})\right]\theta(t-t_{0_{QQ}})$$

$$+2\frac{M_{K^*}^2}{g_{K^*}^2}\,\frac{M_Q^2}{g_Q^2}\sqrt{(t-t_{1_{K^*Q}})(t-t_{0_{K^*Q}})}\left[3(t_0-t_1)_{K^*Q}^2+4(t-t_{0_{K^*Q}})(t-t_{1_{K^*Q}})\right]$$

$$\theta(t-t_{0_{K^*Q}})$$

$$+16f_K^2\,\frac{M_{K^*}^2}{g_{K^*}^2}\left[(t-t_{1_{K^*K}})(t-t_{0_{K^*K}})\right]^{3/2}\theta(t-t_{0_{K^*K}})$$

... / ...

$$+ 16 f_K^2 \frac{M_Q^2}{g_Q^2} \left[(t - t_{1KQ})(t - t_{0KQ}) \right]^{3/2} \theta(t - t_{0KQ})$$

$$+ 4 f_K^4 \sqrt{t(t - t_{0KK})} \left[t_{0KK}^2 + tt(t - t_{0KK}) \right] \theta(t - t_{0KK}) \Big\} . \qquad (14)$$

The notations are:

$$t_{0KQ} = (M_K + M_Q)^2 \ , \quad t_{1KQ} = (M_Q - M_K)^2 \ \text{etc...} \ . \qquad (15)$$

The QCD continuum can be computed from Figs. (6a,b)[4] and is:

$$\text{Im}\ \psi(t) = \frac{1}{16} \left(\frac{1}{16\pi^2} \right)^3 \frac{8}{5} \left(1 + 1/3 \right) t^4 \ \theta(t - t_c) \ . \qquad (16)$$

We make it start from a threshold t_c high enough to avoid double counting with the low energy hadronic contributions.

Comparing Eqs. (14)-(16) with the inequality Eq. (10) and taking the limit $t \to t_{0KK} + \varepsilon$, we obtain the bound:

$$| F(t_{0KK}) |^2 \leq \frac{1}{16} (1 + 1/3) \left(4 f_K^2 M_K^2 \right)^2 + \mathcal{O}\left(\frac{1}{N_c^2} \right) . \qquad (17)$$

The bound (17), obtained in the narrow width approximation, only depends on the kaon contribution, not on higher resonances or the continuum. This is a consequence of taking the limit $t \to t_{0KK} + \varepsilon$. It specially makes unnecessary the replacement of the narrow width approximation for the K^* by a more refined evaluation of the $K\pi$ continuum, since the associated hadronic threshold in Im $\psi(t)$ would anyhow be higher than t_{0KK}. The same reasoning applies to other resonances and to the QCD continuum.[F1]

To obtain abound for $|\mathcal{M}_G|^2$, we must complete the inequality

Eq. (17) by an estimate of

$$\left| \frac{F(t_{0KK})}{F(0)} \right|^2 \equiv \left| \frac{F(t_{0KK})}{\mathcal{M}_G} \right|^2 . \qquad (18)$$

The simplest possibility is to take this ratio equal to 1. The KK, I = 1, J = 0, is indeed an exotic channel most probably structureless (though we have still no experimental result on it). A slightly more refined argument can be given in the line of Ref. [4], by linking this channel to the other exotic one, ππ I = 2, J = 0, belonging to the same representation (27) of flavor SU(3). This last one has been experimentally studied at low energy. A good fit to the data for the phase shift δ_0^2 is given by the effective range formula:

$$\delta_0^2 = \text{Arctg} \left[\frac{1}{\frac{a_0^2}{2} \sqrt{(t - t_{0\pi\pi})}} + \frac{1}{2} r_0^2 \frac{1}{2} \sqrt{(t - t_{0\pi\pi})} \right]^{-1} , \qquad (19)$$

with:[7]

$$a_0^2 \simeq -.15\ f \simeq -.75\ \text{GeV}^{-1} \ ,$$

$$r_0^2 \simeq .13\ f \simeq .65\ \text{GeV}^{-1} \ . \qquad (20)$$

δ_0^2 is related to F(t)/F(0) by an Omnes relation[8]

$$\frac{F(t)}{F(0)} = P_n(t) \ \exp \left[\frac{t}{\pi} \oint_{t_0}^{\infty} \frac{dt'}{t'} \ \frac{\delta_0^2(t')}{t'(t' - t)} \right] . \qquad (21)$$

$P_n(t)$ is a polynomial of n^{th} degree in t, n being the number of zeros of F(t) in the complex plane. (Note that we neglect all effects of inelasticity.) Taking the simplest assumption of the absence of zero,[F2] and so $P_n(t) = 1$, we obtain the estimate

$$\left| \frac{F(t_{0\pi\pi})}{F(0)} \right|_{\pi\pi,\, I=2,\, J=0} \simeq .9 \quad , \tag{22}$$

which we shall take as an estimate of the corresponding quantity in the KK channel $I = 1, J = 0$.

This leads to the final bound:

$$|\mathcal{M}| = |F(0)| \leq 1.1 \sqrt{1+1/3}\, F_K^2 M_K^2 = \frac{1.1}{\sqrt{1+1/3}} |\mathcal{M}|_{vac} = .95 |\mathcal{M}|_{vac} . \tag{23}$$

where $|\mathcal{M}|_{vac}$ is given by Eq.(4). We recall that the sources of uncertainty are:

-the contributions of order $1/N_c^2$,

-PCAC for the kaon, Eq. (11c),

-the estimation of $|F(t_0)/F(0)|_{KK\, I=1,\, J=0}$ from the similar in the $\pi\pi$, $I=2, J=0$ channel and the corresponding low energy elastic phase shift, together with the hypothesis of the absence of zero for F(t). The second uncertainty, essentially attached to the existence of a continuum higher in energy in addition to the kaon pole, is again of no relevance in the limit $t \to t_{0KK} + \varepsilon$. The third is certainly the more out of control. We are however inclined to trust the intuitive idea that in a smooth exotic channel, the scalar form factor cannot vary very much within the small domain of energy $[0, 4M_K^2]$.

Regarding to that and to Eq. (22) we shall attach an uncertainty of ordr 10% to our last following statement: Up to $1/N_c^2$ corrections, the vacuum estimate is an upper bound for the matrix element of the local

four quark operator O(x) between the states K^0 and $\bar{K}^0$. This strengthens the upper bound of $2|\mathcal{M}|_{vac}$ obtained previously in Ref.[4] by hadronic sum rules, is consistent with recent Monte-Carlo simulations on a lattice[9] giving $|\mathcal{M}|_{vac}$ as a reliable estimate, and compatible with Ref. [10], giving $|\mathcal{M}| = .33\, |\mathcal{M}|_{vac}$.

III. THE LAPLACE SUM RULE APPROACH

Both sides of the inequality (10), positive functions of t, may be integral transformed with a positive weight function of t, giving a sum rule in its usual form. (A Hilbert transform was for example used in Ref. [4]). It is clear that this technique, mixing all values of t, generally looses information and gives further uncertainty attached to the contributions of higher resonances and the existence of an arbitrary scale q^2(or M^2). We shall show however, for completeness, that we can recover the bound of Section II by using a Laplace sum rule at the limit of small M.

Taking the weight function in the integral to be e^{-t/M^2}, neglecting here color factors and making the reasonable assumption $|F(t)/F(0)| = 1$, we obtain the bound

$$\left|\frac{\mathcal{M}}{f_K^2 M_K^2}\right| \leq \sqrt{2}\, 4\pi \frac{M^4}{f_K^2 M_K^2} \left[\frac{-M^2 \left(\hat{L}\,\psi^{(5)}\right)(M^2)}{\xi(M^2)}\right]^{1/2}, \tag{24}$$

where $\hat{L}$ is the Laplace inverse operator[5]

$$\hat{L} = \lim_{\substack{Q^2 \to \infty \\ n \to \infty}} \; {}_{Q^2/n = M^2 \text{ fixed}} \; \frac{(-1)^n}{(n-1)!} (Q^2)^n \frac{\partial^n}{(\partial Q^2)^n} . \tag{25}$$

$\psi^{(5)}(Q^2)$ is the fifth derivative of $\psi(Q^2)$ with respect to Q^2 (as requested by the QCD computation of Fig. (6);[4] we have the identity

$$-M^2\left(\hat{L}\,\psi^{(5)}\right)(M^2) = \frac{1}{(M^2)^5} \frac{1}{\pi} \int_0^\infty dt\; e^{-t/M^2}\, \mathrm{Im}\,\psi(t) \quad . (26)$$

$$\xi(M^2) = \frac{1}{M^2} \int_{t_{0KK}}^{\infty} dt\; e^{-t/M^2} \sqrt{\frac{t - t_{0KK}}{t}} \quad . \tag{27}$$

The different contributions to $-M^2 L\psi^{(5)}$ are plotted on Fig. (7), showing clearly the damping of high energy contributions by the exponential weight at low M. The corresponding bound for $|\mathcal{M}/f_K^2 M_K^2|$ is shown on Fig. (8); we recover for $M \to 0$, as expected, the previous bound (up to a color factor and the variation of $F(t_0)$ to $F(0)$), which constitutes in this case the maximum capability of the sum rule technique.

IV. CONCLUSION

We have obtained the upper bound:

$$\left| \langle \bar{K}^0 | \left(\bar{s}_L \gamma^\mu d_L \; \bar{s}_L \gamma_\mu d_L \right)_{local} | K^0 \rangle \right| \leq$$

$$(\text{vacuum estimate}) + \mathcal{O}\left(1/N_c^2\right),$$

with an estimated uncertainty of order 10%.

ACKNOWLEDGMENTS

The author wants to acknowledge the hospitality of the theory group at Lawrence Berkeley Laboratory and correspondance with B. Guberina. This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

REFERENCES

[1.] M. K. Gaillard, B. W. Lee, *Phys. Rev.* D10, 897 (1974).

[2.] M. I. Vysotskii, *Sov. J. Nucl. Phys.* 31, 797 (1980).

[3.] A. I. Vainshtein, V. I. Zakharov, V. A. Novikov, M. A. Shifman, *Sov. J. Nucl. Phys.* 23, 540 (1976).
F. J. Gilman, M. B. Wise, *Phys. Lett.* B83, 83 (1979);
B. Gaiser, T. Tsao, M. B. Wise, *Ann. Phys.* 132, 66 (1981);
F. J. Gilman, M. B. Wise, *Phys. Rev.* D27, 1128 (1983).

[4.] B. Guberina, B. Machet, E. de Rafael, *Phys. Lett.* 128B, 269 (1983).

[5.] M. A. Shifman, A. I. Vainshtein, V. I . Zakharov, *Nucl. Phys.* B147, 385 (1979).
S. Narison, E. de Rafael, *Phys. Lett.* 103B, 57, (1981);
D. V. Widder "The Laplace Transform": Princeton University Press (1946).

[6.] G. 'tHooft, *Nucl. Phys.* B75, 461 (1974).

[7.] D. Cohen, T. Ferbal, P. Slattery, B. Werner, *Phys. Rev.* D7, 661 (1973).

[8.] (a).R. Omnes, *Nuov, Cim.* 8, 316 (1958).
(b).B. R. Martin, D. Morgan, G. Shaw:
"Pion-pion interaction in particle physics" p 176 Academic Press (1976).

[9.] R. C. Brower, G. Maturana, M. B. Gavela, R. Gupta, Harvard preprint HUTP 84/A004 (1984).
N. Cabibbo, G. Martinelli, R. Petronzio, preprint TH 3774-CERN (1983).

[10]. J. F. Donoghue, E. Golowich, B. R. Holstein, *Phys. Lett.* 119B, 412 (1982).

FOOTNOTES

F1. It is worth mentioning that, even in the sum rule approach of Ref. [4], mixing all values of t, the bounds obtained with a narrow width approximation for the K* or the more elaborate form of the Kπ continuum used therein do not differ by more than 6% (at the minimum of the curve c of Fig. 3.)

F2. Up to now this hypothesis is in contradiction with no experimental result. A possible test would be the violation of the sum rule[8b]

$$0 = \int_{4m_\pi^2}^{\infty} dt' \, \frac{\ln |F_\pi(t')|}{t'(t'-t_{0\pi\pi})^{1/2}}$$

very difficult to detect experimentally.

FIGURE CAPTION

Fig. 1 K^0K^0 transition by exchanges of W's.

Fig. 2. Contributions that we neglect in the local approximation.

Fig. 3. The propagator $\psi(q^2)$.

Fig. 4. Factorization of $\psi(q^2)$.

Fig. 5. Factorization only breaks down at order $1/N_c^2$.

Fig. 6. Free quark contributions to $\psi(q^2)$.

Fig. 7. Contributions of K, K*, Q to $-M^2[L\,\psi^{(5)}]\,(M^2)$.

Fig. 8. Bound for $|\mathcal{M}/f_K^2M_K^2|$ in the Laplace sum rule technique.

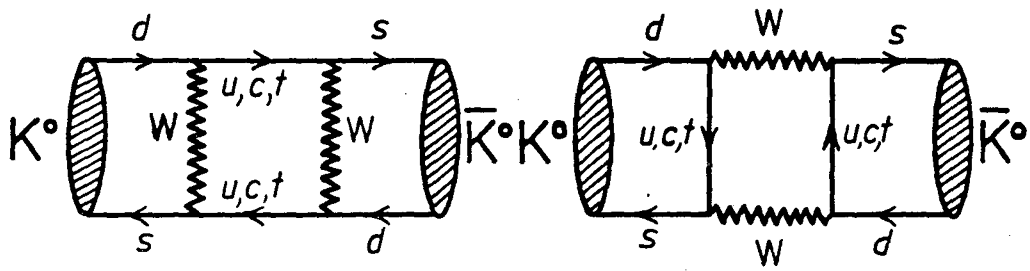


Figure 1

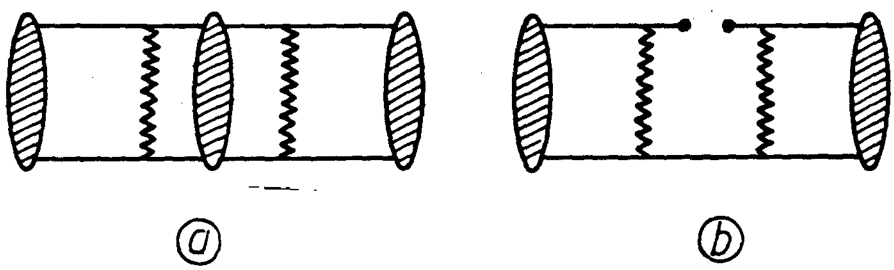


Figure 2

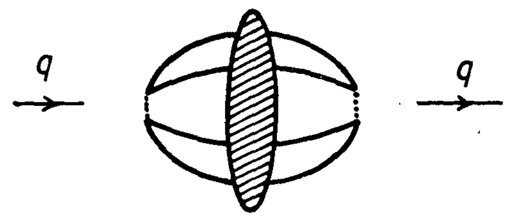


Figure 3

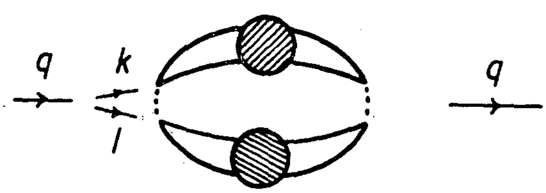


Figure 4

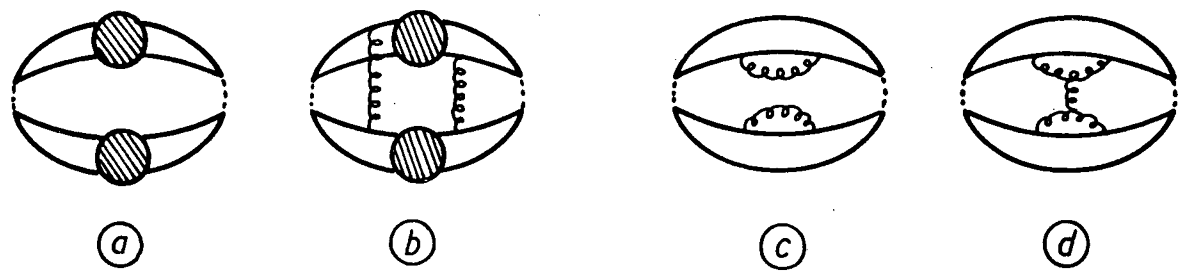


Figure 5

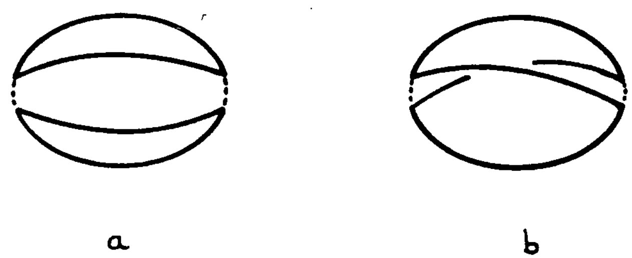
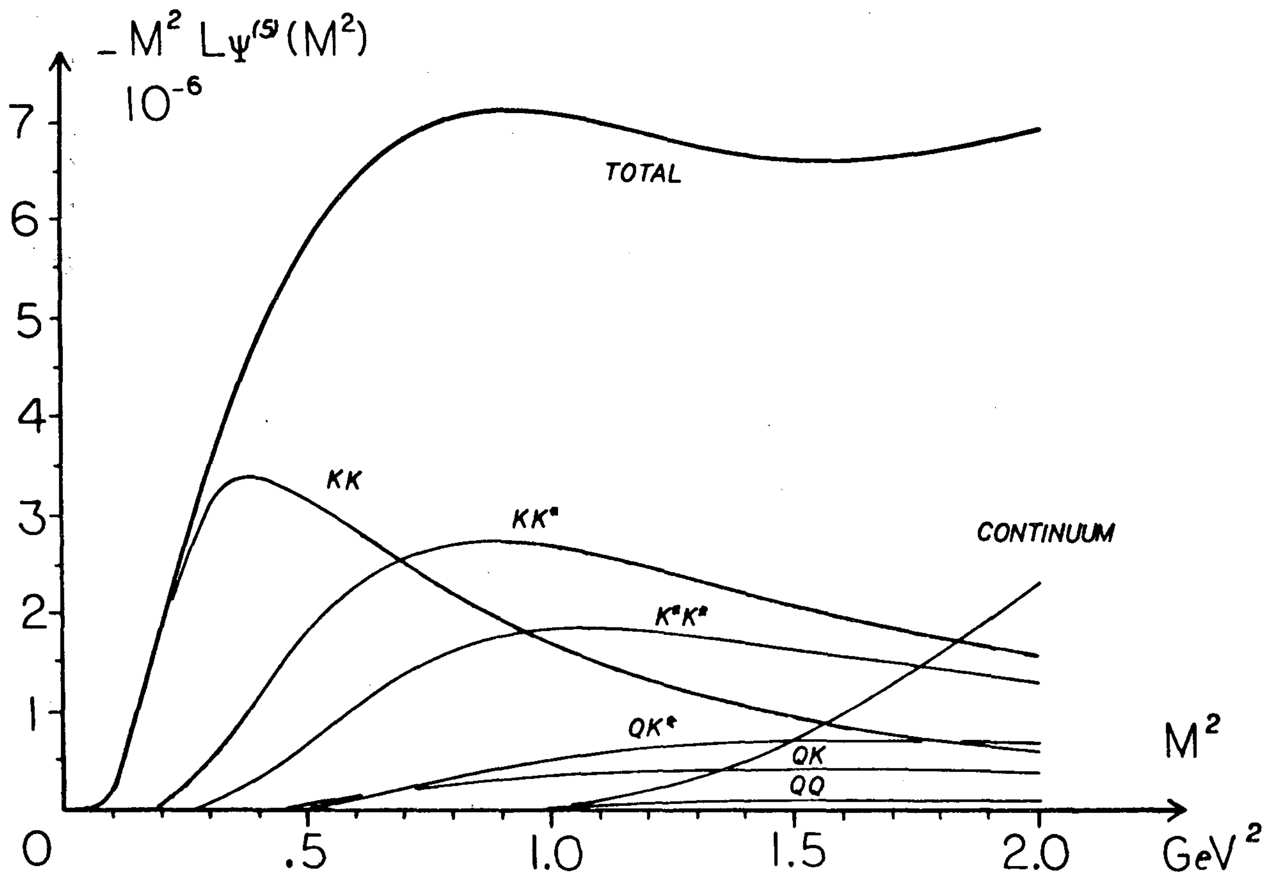


Figure 6



$-M^2 L\psi^{(5)}(M^2)$
10^{-6}
TOTAL
KK
KK*
K*K*
CONTINUUM
QK*
QK
QQ
M^2
GeV²
0
.5
1.0
1.5
2.0
1
2
3
4
5
6
7

Figure 7

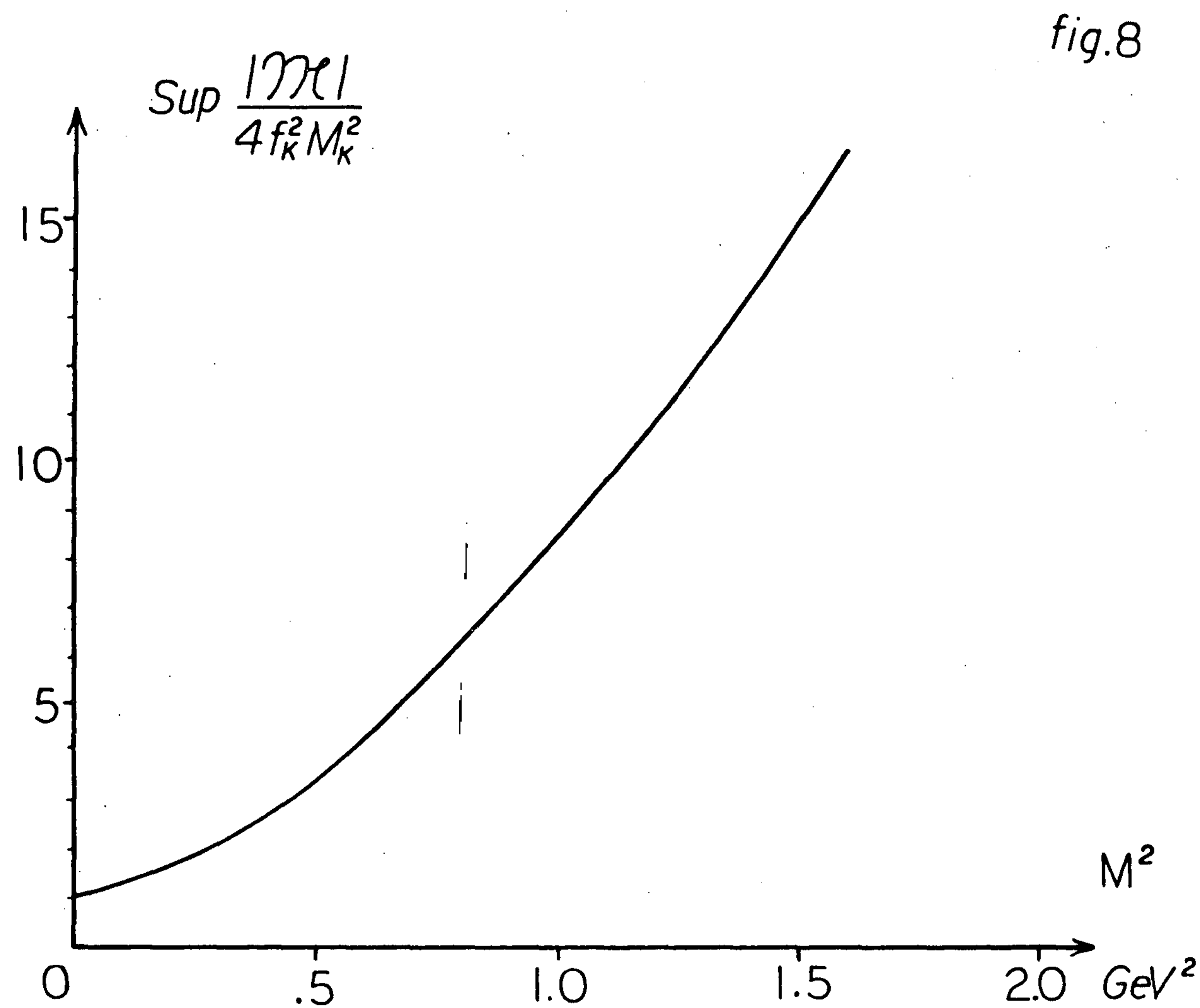


Figure 8

TECHNICAL INFORMATION DEPARTMENT
LAWRENCE BERKELEY LABORATORY
UNIVERSITY OF CALIFORNIA
BERKELEY, CALIFORNIA 94720