CM-P00049206

Ref.TH.1734-CERN

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DUALITY IN THE CURRENT PROPAGATOR

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A B S T R A C T

The idea of local duality between generalized vector dominance and light cone algebra or parton model is applied to the propagator of vector currents. A set of sum rules reminiscent of the generalized first Weinberg sum rules but involving certain unknown parameters are derived. These parameters are fixed from the leptonic decay modes of ρ , ω , ϕ . The sum rules are used to calculate the $\omega - \phi$ mixing angle and thereby predict $\Gamma(\phi \rightarrow K\bar{K}) = 3.25 \pm 0.34$ MeV, in agreement with experiment. Simple mass mixing or current mixing model for the vector mesons seem to be inconsistent with data. The total cross-section for $e^+e^- \rightarrow$ hadrons is calculated in the scaling region. If scaling is achieved at $s \simeq 9-16 \text{ GeV}^2$ then it is argued that about 22% of the events in this energy region will include strange particles.

1. - INTRODUCTION

In a recent series of papers Greco et al. ¹⁾ and Sakurai and Schildknecht ²⁾ have indicated that the pointlike behaviour of the interactions of massive photons can be assimilated by assuming the existence of an infinite series of vector states coupled to this photon. This assumption was called generalized vector meson dominance (GVMD). Subsequently, Sakurai ³⁾ and Böhm, Joos and Kramer ⁴⁾ have conjectured that there is a duality between vector meson formation and pointlike $1/s$ behaviour for the total cross-section $\sigma(e^+e^- \rightarrow \text{hadrons})$. This can be thought to imply that for the absorptive part of the current propagator the GVMD description and the light cone algebra (LCA) or parton model (PM) description are just dual to each other. Here we take this point of view. We can then use finite energy sum rules (FESR) to calculate the high s behaviour from the low s data in a way analogous to the one applied to strong interaction process ³⁾. The basic assumption of this paper is that the lowest vector mesons are enough to saturate these sum rules ³⁾. We term this as local duality (LD). This assumption presumably applies to all cases where an off-mass shell photon is involved. The case of the current propagator is the easiest to handle, since no special assumptions about the couplings of the various vector mesons to an hadronic system is necessary.

LD leads to a set of sum rules reminiscent of the generalized first Weinberg sum rules ⁵⁾. Using them and a reasonable assumption on the φ NN coupling, we fix the ω - φ mixing angle and predict the decay rate $\Gamma(\varphi \rightarrow K\bar{K})$ in terms of $\Gamma(\varphi \rightarrow e^+e^-)$. The prediction is in very good agreement with experiment. This is done in Sections 2 and 3, where some other consequences of our calculations are also presented. In Section 4 the total cross-section for $e^+e^- \rightarrow \text{hadrons}$ in the scaling region is calculated. Finally, in Section 5 we summarize our conclusions.

2. - THE PROPAGATOR OF THE VECTOR CURRENTS

The absorptive part of the current propagator is given by

$$\begin{aligned} \frac{1}{2} \int d^4x e^{-iqx} \langle 0 | j_\mu^\alpha(x) j_\nu^\beta(0) | 0 \rangle = \\ = \pi (q_\mu q_\nu - q^2 \delta_{\mu\nu}) R^{\alpha\beta}(s) + q_\mu q_\nu R_1^{\alpha\beta}(s). \end{aligned} \quad (1)$$

Disregarding the non-conserved parts of the currents, we obtain

$$R^{\alpha\beta}(s) = \frac{(2\pi)^3}{3s} \sum_n \delta^{(4)}(p_n - q) \langle 0 | j_\mu^\alpha(0) | n \rangle \langle n | j_\nu^\beta(0) | 0 \rangle. \quad (2)$$

Here $s = -q^2$ and $\alpha, \beta = 0, \dots, 8$ are the U_3 indices. All the results derived here refer to the conserved currents corresponding to $\alpha, \beta = 0, 3, 8$.

Parton model (PM) or light cone algebra (LCA) or simply asymptotic U_3 imply that ⁶⁾

$$\lim_{s \rightarrow \infty} R^{\alpha\beta}(s) = \delta_{\alpha\beta} C, \quad (3)$$

where C is a constant whose value seems to depend on the particular version of PM or LCA used ⁶⁾. According to the philosophy of the present work, Eq. (3) can also be realized assuming the existence of an infinite series or a continuum of vector states coupled to the currents ^{1), 2)}. For the calculations presented here, it is convenient to express $R^{\alpha\beta}(s)$ in terms of spectral functions in which at most one of the lowest vector mesons ρ, ω, φ contributes. To this end we introduce new currents $\tilde{j}_\mu^\alpha, \tilde{\omega}_\mu$ and $\tilde{\varphi}_\mu$ and related spectral functions as follows

$$\begin{aligned} j_\mu^\alpha &= \frac{m_\rho^2}{f_\rho} \tilde{j}_\mu^\alpha \quad \alpha = 1, 2, 3, \\ j_\mu^8 &= \frac{2}{\sqrt{3}} j_\mu^0 = \frac{1}{f_\gamma} \left\{ \cos \Theta_Y m_\varphi^2 \tilde{\varphi}_\mu + \sin \Theta_Y m_\omega^2 \tilde{\omega}_\mu \right\}, \\ j_\mu^0 &= \sqrt{\frac{2}{3}} j_\mu^8 = \frac{1}{f_\rho} \left\{ -\sin \Theta_B m_\varphi^2 \tilde{\varphi}_\mu + \cos \Theta_B m_\omega^2 \tilde{\omega}_\mu \right\}, \end{aligned} \quad (4)$$

and *)

$$\begin{aligned}
 R_g(s) &= \frac{(2\pi)^3}{3s} m_g^4 \sum_n \delta^{(4)}(p_n - q) \langle 0 | \tilde{g}_r | n \rangle \langle n | \tilde{g}_r | 0 \rangle , \\
 R_w(s) &= \frac{(2\pi)^3}{3s} m_w^4 \sum_n \delta^{(4)}(p_n - q) \langle 0 | \tilde{w}_r | n \rangle \langle n | \tilde{w}_r | 0 \rangle , \\
 R_y(s) &= \frac{(2\pi)^3}{3s} m_y^4 \sum_n \delta^{(4)}(p_n - q) \langle 0 | \tilde{y}_r | n \rangle \langle n | \tilde{y}_r | 0 \rangle , \\
 R_{yw}(s) &= \frac{(2\pi)^3}{3s} m_y^2 m_w^2 \sum_n \delta^{(4)}(p_n - q) \langle 0 | \tilde{y}_r | n \rangle \langle n | \tilde{w}_r | 0 \rangle .
 \end{aligned} \tag{5}$$

They are defined such that

$$\langle 0 | \tilde{g}_r(0) | g \rangle = \langle 0 | \tilde{w}_r(0) | w \rangle = \langle 0 | \tilde{y}_r(0) | y \rangle = 1 \epsilon_r , \tag{6}$$

$$\langle 0 | \tilde{y}_r | w \rangle = \langle 0 | \tilde{w}_r | y \rangle = 0 \tag{7}$$

where ϵ_r is the spin wave function of the relevant vector meson.

From (4) and (5)

$$\begin{aligned}
 R^{33}(s) &= \frac{1}{f_r^2} R_g(s) , \\
 R^{88}(s) &= \frac{3}{4} \frac{1}{f_y^2} \left\{ \cos^2 \theta_y R_y(s) + \sin^2 \theta_y R_w(s) + \sin(2\theta_y) R_{yw}(s) \right\} , \\
 R^{00}(s) &= \frac{3}{2} \frac{1}{f_B^2} \left\{ \sin^2 \theta_B R_y(s) + \cos^2 \theta_B R_w(s) - \sin(2\theta_B) R_{yw}(s) \right\} , \\
 R^{80}(s) &= \frac{3}{2\sqrt{2}} \frac{1}{f_y f_B} \left\{ -\cos \theta_y \sin \theta_B R_y(s) + \sin \theta_y \cos \theta_B R_w(s) + \right. \\
 &\quad \left. + \cos(\theta_y + \theta_B) R_{yw}(s) \right\} .
 \end{aligned} \tag{8}$$

*) $R_{yw}(s)$ is real because of time inversion invariance.

Equations (3), (8) imply that each of the spectral functions defined in (5) approaches some constant value at high energies. We relate these constant values to the properties of the low vector mesons by using the assumption of local duality³⁾. The precise meaning of this assumption is that we can write FESR for each of the spectral functions and saturate them by only the lowest vector mesons ρ , ω , φ . Let us apply this technique to $R_\rho(s)$. By assuming

$$\lim_{s \rightarrow \infty} R_\rho(s) = C_\rho \quad (9)$$

we write the FESR

$$\int_{s_0}^{a_\rho m_\rho^2 + s_0} R_\rho(s) ds = C_\rho a_\rho m_\rho^2 \quad (10)$$

The quantity $a_\rho m_\rho^2$ defines the region around the ρ pole over which the integral has to be calculated so that (9) and (10) are satisfied. In a rough sense, $a_\rho m_\rho^2$ describes the average distance (in mass square) of successive vector mesons ρ , ρ' , ρ'' , etc. Its inverse can be thought of as the density of the ρ -like states. We find⁷⁾

$$\lim_{s \rightarrow \infty} R_\rho(s) = C_\rho = \frac{1}{a_\rho} \quad (11a)$$

Using similar relations for the other spectral functions and defining the quantities $a_\omega m_\omega^2$, $a_\varphi m_\varphi^2$ in an analogous way, we obtain^{*)}

$$\lim_{s \rightarrow \infty} R_\omega(s) = C_\omega = 1/a_\omega \quad , \quad (11b)$$

$$\lim_{s \rightarrow \infty} R_\varphi(s) = C_\varphi = 1/a_\varphi \quad , \quad (11c)$$

$$\lim_{s \rightarrow \infty} R_{\varphi\omega}(s) = 0 \quad . \quad (11d)$$

*) In the regions where the notion of ω - and φ -like states makes sense, we can consider $1/(a_\omega m_\omega^2)$, $1/(a_\varphi m_\varphi^2)$ as the respective average densities for these states (see below).

Equations (11a-c) are not very interesting since c_s , a_s , as well as the corresponding quantities for ω , φ , are unknown. The most interesting result is (11d). It follows from the fact that none of the ω , φ or ρ contribute to $R_{\varphi\omega}(s)$. Therefore, it is obtained if we write an FESR for $R_{\varphi\omega}(s)$ and cut it off just after the φ meson. Equation (11d) means that in the continuum (scaling) region, the isoscalar states connected to the photon are coupled mainly either to current $\tilde{\varphi}_\mu$ or $\tilde{\omega}_\mu$. We term these states φ -like and ω -like, respectively. If scaling is achieved for $s \gtrsim 1.5 \text{ GeV}^2$, then all isoscalar states are either φ -like or

ω -like. On the basis of (4), we see that φ -like states should have a quark structure similar to that of the ordinary φ ; i.e., they are mainly $\lambda\bar{\lambda}$ excitations, in the constituent quark frame. Similarly the ω -like states are made out of non-strange quarks. The decay products of a φ -like state will almost always contain strange particles. On the contrary there will be a considerable probability that no strange particles are produced in a decay of an ω -like state, provided that its mass is not very high *).

From (3), (8), (11) we get

$$(a_\varphi/a_\omega)^{1/2} \tan \theta_Y = (a_\omega/a_\varphi)^{1/2} \tan \theta_B, \quad (12a)$$

$$\begin{aligned} \frac{1}{f_\rho^2} \frac{1}{a_\varphi} &= \frac{3}{4} \frac{1}{f_Y^2} \left\{ \frac{\cos^2 \theta_Y}{a_\varphi} + \frac{\sin^2 \theta_Y}{a_\omega} \right\} = \\ &= \frac{3}{2} \frac{1}{f_B^2} \left\{ \frac{\sin^2 \theta_B}{a_\varphi} + \frac{\cos^2 \theta_B}{a_\omega} \right\}. \end{aligned} \quad (12b)$$

It is convenient to define

$$\tan^2 \theta \equiv \tan \theta_Y \tan \theta_B \quad (13)$$

*) This is based on the empirical fact that very few events involving strange particles are observed in $p\bar{p}$ or pp scattering at $s = 9-16 \text{ GeV}^2$ (see Section 4).

Solving (12) and substituting in (4) we find

$$\begin{aligned}
 J_\mu^3(x) &= (m_\rho^2 / t_\rho) \tilde{J}_\mu^3(x) \\
 J_\mu^Y &= \frac{2}{\sqrt{3}} J_\mu^P = \frac{2}{\sqrt{3} f_\rho} \left\{ \cos \theta \left(\frac{a_\rho}{a_\rho} \right)^{1/2} m_\rho^2 \tilde{J}_\mu^P + \sin \theta \left(\frac{a_\omega}{a_\rho} \right)^{1/2} m_\omega^2 \tilde{J}_\mu^P \right\}, \\
 J_\mu^0 &= \sqrt{\frac{2}{3}} J_\mu^0 = \sqrt{\frac{2}{3}} \frac{1}{f_\rho} \left\{ -\sin \theta \left(\frac{a_\rho}{a_\rho} \right)^{1/2} m_\rho^2 \tilde{J}_\mu^P + \cos \theta \left(\frac{a_\omega}{a_\rho} \right)^{1/2} m_\omega^2 \tilde{J}_\mu^P \right\}. \quad (14)
 \end{aligned}$$

Notice that for the derivation of (12) and (14) we essentially need local duality only for $R_{\rho\omega}(s)$. To see this explicitly just replace in these equations

$$(a_\rho, a_\omega, a_\rho) \rightarrow (1/c_\rho, 1/c_\omega, 1/c_\rho), \quad (15)$$

and remember (11a-c).

Equations (14) can be used to fix the mixing angle θ . According to our preceding remarks on the ρ -like states and the known properties of lowest ρ meson

$$\langle N | \tilde{J}_\mu^3(0) | N \rangle_{q^2=0} = 0 \quad (16)$$

Equation (16) can also be argued for on the basis of (ordinary) vector meson dominance for the isoscalar Dirac form factor $F_1^S(t)$ of the nucleon. From (16) and (14), we find that the mixing angle is ideal, i.e.,

$$\tan \theta = \frac{1}{\sqrt{2}} \quad (17)$$

It is important to realize that the derivation of (17) does not use any specific model for the mixing of the $\omega - \rho$ states ^{8),9)}. It is just a consequence of (16), (3) and the local duality [i.e., Eq. (11d)]. To see this, let us drop LD. Then the spectral functions $R_\rho(s), \dots, R_{\rho\omega}(s)$ go asymptotically to unknown constants defined by the left-hand side of (11a-c) and

$$\lim_{s \rightarrow \infty} R_{\varphi\omega}(s) = C_{\varphi\omega}.$$

Substituting (8) in (3) and using (16), and (13), we find

$$\begin{aligned} \left(\frac{1}{2} - t g^4 \theta\right) C_{\omega} + \frac{C_{\varphi\omega}}{2 C_{\varphi}} \left(\frac{\cot^2 \theta}{2} + t g^2 \theta \right) \left[C_{\varphi\omega} (\cot \theta - t g \theta) + \right. \\ \left. + \sqrt{C_{\varphi\omega}^2 (\cot \theta - t g \theta)^2 + 4 C_{\varphi} C_{\omega}} \right] = 0 \end{aligned} \quad (18)$$

We see that (18) gives (17) only if the LD result $C_{\varphi\omega} = 0$ is used ^{*)}.

We can also turn the statement around and say that (17) implies the LD result (11d). In the next section, we use (17) to predict $\Gamma(\varphi \rightarrow K\bar{K})$.

Let us define

$$f_{\omega} = \left(\sqrt{3} f_{\rho} / \sin \theta \right) \left(a_{\rho} / a_{\omega} \right)^{1/2} = 3 f_{\rho} \left(a_{\rho} / a_{\omega} \right)^{1/2}, \quad (19a)$$

$$f_{\varphi} = \left(\sqrt{3} f_{\rho} / \cos \theta \right) \left(a_{\rho} / a_{\varphi} \right)^{1/2} = \frac{3 f_{\rho}}{\sqrt{2}} \left(a_{\rho} / a_{\varphi} \right)^{1/2}. \quad (19b)$$

Parameters f_{ρ} , f_{ω} , f_{φ} are determined from the leptonic decay modes of ρ , ω , φ respectively. Using the experimental results ¹⁰⁾, we find

$$\begin{aligned} \frac{a_{\omega} m_{\omega}^2}{a_{\rho} m_{\rho}^2} = \frac{9 f_{\rho}^2 m_{\omega}^2}{f_{\omega}^2 m_{\rho}^2} = 1.30 \pm 0.17 \\ \frac{a_{\varphi} m_{\varphi}^2}{a_{\rho} m_{\rho}^2} = \frac{9 f_{\rho}^2 m_{\varphi}^2}{2 f_{\varphi}^2 m_{\rho}^2} = 1.50 \pm 0.23 \end{aligned} \quad (20)$$

Both these numbers are not very far from unity. This is encouraging for the validity of LD for the spectral functions $R_{\rho}(s)$, $R_{\varphi}(s)$, $R_{\omega}(s)$. Indeed, Eqs. (20) mean that the average densities of the φ - and ω -like states are about the same and not much different from the density of the ρ -like states. Such a result is reasonably close to the strict Veneziano spectrum which demands common densities for ρ , ω - and φ -like states; namely,

^{*)} Notice from (5) that $C_{\varphi} C_{\omega} \geq C_{\varphi\omega}^2$.

$$\alpha_\rho m_\rho^2 = \alpha_\omega m_\omega^2 = \alpha_\gamma m_\gamma^2. \quad (21)$$

It is amusing to remark that substituting (21) in (12) we obtain the relations previously derived by saturating the generalized first Weinberg sum rules by the lowest vector mesons ⁵⁾. The relation thus obtained from the right-hand part of (12b) can be directly tested by the experiment. It is found to be violated by more than three standard deviations ¹⁰⁾. In our theory, (12) is satisfied by construction and, instead of (21), we have (20).

We now present the implications of our calculations on the mixing model of vector mesons. Using (12a), (20)

$$\tan \theta = \left(f_\omega / (\sqrt{2} f_\gamma) \right) \tan \theta_\gamma = \left(\sqrt{2} f_\gamma / f_\omega \right) \tan \theta_0, \quad (22)$$

from which ¹⁰⁾

$$\begin{aligned} \theta_\gamma &= 45.5^\circ \pm 1.7^\circ, \\ \theta_0 &= 30.6^\circ \pm 1.6^\circ. \end{aligned} \quad (23)$$

Now the mass mixing model for vector mesons demands ⁸⁾

$$\tan \theta_\gamma = \tan \theta_0 \quad (\text{mass mixing}),$$

while the current mixing model ⁸⁾

$$\frac{m_\omega^2}{m_\gamma^2} \tan \theta_\gamma = \tan \theta_0 \quad (\text{current mixing}).$$

Comparing these relations to our result (22), we find that the mass mixing is violated in our theory by more than three standard deviations since ¹⁰⁾

$$\frac{f_\omega^2}{2 f_\gamma^2} = 0.696 \pm 0.089 \neq 1,$$

while the current mixing is violated by a little more than one standard deviation

$$\frac{f_{\omega}^2}{2f_{\eta}^2} = 0.696 \pm 0.089 \neq \frac{m_{\omega}^2}{m_{\eta}^2} = 0.59$$

We conclude that in the limit where (16) is exact, the LD assumption is inconsistent with either the mass mixing or the current mixing model.

3. - PROCESS $\varphi \rightarrow K\bar{K}$

Let us assume that the matrix elements

$$\langle K^+ | \tilde{g}_r | K^+ \rangle, \langle K^+ | \tilde{\omega}_r | K^+ \rangle, \langle K^+ | \tilde{g}_r | K^+ \rangle \quad (24)$$

obey ordinary VMD for small momentum transfers. This assumption is motivated by the known fact that the pion form factor also seems to obey ^{11), 10)} VMD. Applying it to the first of the above matrix elements and using only the definitions (4), (13), we obtain ⁸⁾

$$\Gamma(\varphi \rightarrow K\bar{K}) \Gamma(\varphi \rightarrow e^+e^-) = \frac{\alpha^2}{18} \frac{p_{KK}^3}{m_{\varphi}} \cos^4 \theta \quad (25)$$

From (17) and $\Gamma(\varphi \rightarrow e^+e^-)$ we predict ⁸⁾

$$\Gamma(\varphi \rightarrow K\bar{K})_{th} = 3.24 \pm 0.34 \text{ MeV} \quad (26)$$

to be compared with the experimental result ¹⁰⁾

$$\Gamma(\varphi \rightarrow K\bar{K})_{Exp} = 3.44 \pm 0.23 \text{ MeV}.$$

The agreement between theory and experiment is very good. We point out that our procedure explains $\Gamma(\varphi \rightarrow K\bar{K})$ without need of any appreciable width effects for the φ meson. This is intuitively appealing since the width is small, but contrary to previous treatments ¹²⁾. This result is also independent of any specific model on the mixing of vector mesons ¹³⁾.

Applying now VMD to all matrix elements in (24) and using (17) we have the K form factor for small momentum transfers

$$F_K(s) = \frac{1}{2} \left\{ F_\pi(s) + \frac{2}{3} \frac{m_\rho^2}{m_\rho^2 - s - i m_\rho \Gamma_\rho} + \frac{1}{3} \frac{m_\omega^2}{m_\omega^2 - s - i m_\omega \Gamma_\omega} \right\} \quad (27)$$

where ¹⁰⁾

$$F_\pi(s) = \frac{m_\rho^2 \left(1 + \delta \frac{\Gamma_\rho}{m_\rho}\right)}{m_\rho^2 - s - i m_\rho \Gamma_\rho \left(k/k_\rho\right)^2 \left(m_\rho/\sqrt{s}\right)} + (s - \omega)_{\text{mixing term}}$$

is the pion form factor ^{*}). The result (27) is the same as the simple SU₃ prediction and is known to agree with the existing experimental information at ¹⁴⁾ $s = 2.4 \text{ GeV}^2$.

4. - PROCESS $e^+e^- \rightarrow \text{HADRONS}$

We now turn to the study of process

$$e^+e^- \rightarrow \text{hadrons} \quad (28)$$

It is well known that

$$\frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\left(\frac{4\pi\alpha^2}{3s}\right)} = 3\pi R^{\text{em}}(s) \quad , \quad (29)$$

where $R^{\text{em}}(s)$ is the spectral function defined in (2) for the electromagnetic current. Using (8), (14), (17), we find for the scaling region

$$\frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \frac{3\pi}{a_\rho} \left\{ \frac{4\pi}{f_\rho^2} + \frac{1}{9} \left(\frac{4\pi}{f_\rho^2}\right) + \frac{2}{9} \left(\frac{4\pi}{f_\rho^2}\right) \right\} \quad (30)$$

^{*}) $\delta = 0.48$, $k_\rho = \frac{1}{2} (m_\rho^2 - 4m_\pi^2)^{1/2}$, $k = \frac{1}{2} (s - 4m_\pi^2)^{1/2}$.

The first term within brackets gives the contribution of φ -like states. The second and third terms come respectively from the ω -like and φ -like states. We see that the isoscalar photon contribution is $1/3$ the contribution of the isovector photon. This is a direct consequence of (3). We now study the probability P_s that an event containing strange particles will appear in (28).

The decay products of a φ -like state almost always include strange particles. We write, in the scaling region,

$$P_s = P_\varphi + (1 - P_\varphi) p_s \quad (31)$$

where

$$P_\varphi = \frac{\left(\frac{2}{9}\right) \frac{4\pi}{f_\rho^2}}{\left(\frac{12}{9}\right) \frac{4\pi}{f_\rho^2}} = \frac{1}{6} \quad (32)$$

is the probability that a φ like state is produced. The parameter p_s describes the probability that the decay products of a φ -like or ω -like state will involve strange particles ^{*}). A fair guess for p_s is to assume that it is the same as the corresponding probability in pp or $p\bar{p}$ scattering. This is motivated on the ground that in both cases the interaction of non-strange quarks is involved. For $s \sim 9-16 \text{ GeV}^2$, we have ¹⁵⁾

$$p_s = (7 \pm 3) \% \quad (33)$$

Substituting (33) in (31), we argue that if scaling has been achieved at energies $s \simeq 9.16 \text{ GeV}^2$, then

$$P_s \simeq 0.22 \pm 0.02 \quad (34)$$

at these energies. As energy increases, we expect that both p_s , P_s will increase until they reach an asymptotic value close to unity at infinite energies.

^{*}) For simplicity it is assumed that this probability is the same for φ -like and ω -like states.

In general we predict that in the scaling region

$$P_s > P_g = 17\% . \quad (35)$$

Equation (34) indicates that a considerable percentage of events containing strange particles is expected even at low energies provided that scaling has been achieved early.

Finally, (30) gives $\sigma(e^+e^- \rightarrow \text{hadrons})$ provided we know a_g . As we mentioned before, $a_g m_g^2$ describes the average distance of successive g -like poles. A g -like state has been seen in Frascati ^{16),17)} at 1.6 GeV. If this is the first such state, it will imply $a_g \simeq 4$. On the other hand there exists also some experimental indication for a g' at the neighbourhood of ¹⁸⁾ B(1235). If this is finally confirmed, it would mean $a_g \simeq 2$ and that the $g(1600)$ should be identified with the ^{1),3)} g'' . Our predictions in these two cases are

$$\frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \begin{cases} 2.5 \pm 0.3 & \text{if } a_g = 2 \\ 1.25 \pm 0.15 & \text{if } a_g = 4 \end{cases} \quad (36)$$

The experimental situation has not been settled yet ^{19),17)}. Nevertheless this result does not seem to be excluded by the experimental data ³⁾. Notice that the result for $a_g = 2$ is in reasonable agreement with Gell-Mann's result for coloured quarks ⁶⁾.

5. - CONCLUSIONS

In this work we applied the concept of local duality to the propagator of the vector currents. The main result is Eq. (11d). Combining this to PM or LCA, the following conclusions are drawn.

- (i) In the scaling region it makes sense to classify the isoscalar states coupled to the currents as φ -like and ω -like. The decay products of a φ -like state almost always contain strange particles.
- (ii) The $\omega - \varphi$ mixing angle is calculated to be ideal. This correctly predicts $\Gamma(\varphi \rightarrow K\bar{K})$ in terms of $\Gamma(\varphi \rightarrow e^+e^-)$ and leads to the conclusion that the K form factor is the same as in the simple SU_3 model.
- (iii) In the limit $g_{\varphi NN} = 0$, the mass mixing model for vector mesons is excluded by three standard deviations while the current mixing model is disfavoured by more than one standard deviation.
- (iv) If scaling for process $e^+e^- \rightarrow \text{hadrons}$ has been achieved at $s \simeq 9-16 \text{ GeV}^2$, then we argue that 22% of the events at this energy region will contain strange particles.

ACKNOWLEDGEMENTS

I would like to thank Prof. J.J. Sakurai and Prof. M. Jacob for enlightening discussions. Thanks are also due to Prof. J.S. Bell for helpful comments and a very critical reading of the manuscript.

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