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Tornqvist, Nils A.

Publication Date

1976-06-01

Submitted to Physics Letters B

LBL-5304
Preprint c. \

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THE OKUBO-ZWEIG-IIZUKA RULE NEAR NEW THRESHOLDS *

Nils A. Törnqvist **†

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

June 7, 1976

Abstract

Because of unitarity we expect the breaking of the Okubo-Zweig-Iizuka (OZI) rule to be abruptly increased at the new threshold ($K\bar{K}$, $D\bar{D}$). Cancellations make the rule much better at high energies and below the threshold. Many experimental puzzles can then be understood.

This work was supported in part by the U. S. Energy Research and Development Administration and in part by the Academy of Finland.

** Participating guest, visiting researcher, Lawrence Berkeley Laboratory.

† Permanent address: University of Helsinki, Research Institute for Theoretical Physics, Helsinki, Finland.

In most recent work on the Okubo-Zweig-Iizuka (OZI) [1] rule many authors have argued using different theoretical concepts, such as asymptotic freedom [2], asymptotic planarity [3] or a dual unitarization scheme [4], that the rule improves monotonically with energy.

In all these approaches one has neglected the effects due to the opening of the thresholds which are allowed by the OZI rule. (e.g., $\psi \rightarrow D\bar{D}$) assuming that their effects are small, - or that they are taken into account in a semilocal duality sense.

It is a well known fact that near thresholds unitarity causes effects in other channels. In ref. [5] it was pointed out that such effects can be important in the OZI breaking. In this paper we shall argue that at the first threshold allowed by the OZI rule there is likely to be a dramatic increase in the breaking.

Let us parametrize the experimental suppression due to the OZI rule by a parameter θ defined by

$$\theta^2 = \frac{x^{\text{n.h.}} \Gamma_{\text{exp}}^{\text{tot}}}{\Gamma_{\text{typ}}^{\text{tot}}}, \quad (1)$$

where

$$\Gamma_{\text{typ}}^{\text{tot}} = 0.2\text{M}, \quad (2)$$

and $x^{\text{n.h.}}$ is the branching ratio to normal hadrons (n.h.). Equation 2 is a fairly good empirical formula for a typical normal hadron resonance not too close to its first threshold. For the psions above 3.9 GeV we do not yet know $x^{\text{n.h.}}$ but with $x^{\text{n.h.}} = 1$ as the upper

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bound one finds[†] using the experimental widths:

If these numbers are interpreted as mixing angles they must be considered as order of magnitude estimates only. Other phenomenological formulae involving phase space, threshold factors and/or specific models can give values differing by a factor 10. The relative magnitudes of the numbers are, however, less sensitive to such considerations.

$$\theta^\psi(3.1) \approx 0.009, \theta^\psi(3.7) \approx 0.005, \theta^\psi(\approx 4) < 0.3. \quad (3)$$

For the ϕ eq. (1) gives $\theta^\phi(1.02) \approx 0.06$. This can be compared with the 7% mixing (in amplitude) of nonstrange quarks in ϕ obtained from a mass matrix analysis (using quadratic masses).

We assume like most authors[3-6] that the bare propagators and vertices of planar amplitudes are ideally mixed and satisfy the OZI rule exactly. The physical amplitudes involve a sum over higher order diagrams including twists. All untwisted loops are assumed to be included in the planar approximation in accordance with the framework of the topological expansion.[7] The simplest OZI breaking diagrams are then those shown in fig. 2a-c which are approximately equivalent through semilocal duality.

The physical propagators (or vertices) then deviate from ideal mixing (or OZI couplings), and we can estimate the magnitude of this deviation by estimating the magnitude of the diagram in fig. 2a-c

compared to a planar diagram.

Using unitarity and some Regge and duality properties one can estimate the discontinuities or the imaginary part of the diagram. Then, assuming normal analyticity properties, we can estimate the real part through a dispersion relation.

Provided one has self-consistency, i.e., that the result turns out to be small (so that only the planar amplitudes are large) one has a well defined dynamical scheme for estimating the OZI breaking.

We argue that the most important discontinuity of the diagram is the discontinuity over the loop, because it involves the product of two (large) planar amplitudes. On the other hand other discontinuities over planar parts of fig. 2a are much smaller, because they involve a product of one planar and one OZI forbidden amplitude. Thus the most important cut starts at the new threshold ($K\bar{K}$ or $D\bar{D}$), whereas the lower kinematically allowed cuts due to normal hadrons are suppressed.

The narrow energy region where only one new OZI allowed threshold is open is particularly simple. Here the discontinuity over a twisted loop is (if one neglects $SU(N)$ breaking) of equal magnitude as that of a (large) planar amplitude since the twists can at most change the sign. Consider the example of D wave, $C = -$, $K\bar{K} \rightarrow K\bar{K}$ and $K\bar{K} \rightarrow \pi\rho$ scattering in the energy region $4m_K^2 \leq s \leq (m_K + m_{K^*})^2$. Both processes are OZI allowed with ω resonances in the s channel, but for $K\bar{K} \rightarrow \pi\rho$ the ϕ resonances are suppressed by the rule. Denote by A^ϕ the part of the amplitude which contains the ϕ poles. We

shall use the fact that for a strong interaction amplitude the discontinuity (or the imaginary part) is of the same order as the amplitude itself. Then using factorization and unitarity we can write at

$$\sqrt{s} = m_\phi$$

$$\frac{A(\phi \rightarrow \pi \rho)}{A(\phi \rightarrow K\bar{K})} = \frac{A^\phi(K\bar{K} \rightarrow \pi \rho)}{A^\phi(K\bar{K} \rightarrow K\bar{K})} \approx \frac{\text{Disc } A^\phi(K\bar{K} \rightarrow \pi \rho)}{\text{Disc } A^\phi(K\bar{K} \rightarrow K\bar{K})}$$

$$= \frac{A^\phi(K\bar{K} \rightarrow K\bar{K}) A(K\bar{K} \rightarrow \rho\pi)^* + A^\phi(K\bar{K} \rightarrow \rho\pi) A(\rho\pi \rightarrow \rho\pi)^*}{A^\phi(K\bar{K} \rightarrow K\bar{K}) A(K\bar{K} \rightarrow K\bar{K})^* + A^\phi(K\bar{K} \rightarrow \rho\pi) A(\rho\pi \rightarrow K\bar{K})^*}$$

$$\approx \frac{A(K\bar{K} \rightarrow \rho\pi)}{A(K\bar{K} \rightarrow K\bar{K})} \quad (4)$$

Here we have used the fact that at the ϕ mass $A^\phi(K\bar{K} \rightarrow K\bar{K}) \approx A(K\bar{K} \rightarrow K\bar{K})$ and in the last step we used the OZI rule (which will be justified a posteriori by the self-consistency) by which terms containing $A^\phi(K\bar{K} \rightarrow \rho\pi)$ can be neglected. Since we only have one angular momentum state a factor $A^\phi(K\bar{K} \rightarrow K\bar{K})$ cancels.

This simple rule relates the suppression factor of an OZI forbidden vertex to a ratio between two OZI allowed amplitudes. Physically it means that, although initially only a kaon pair was emitted, the final state interaction between the kaons is so strong that it determines

which particles eventually emerge.

In order to have self-consistency this ratio must be small for some dynamical reason. In this case we know that $K\bar{K} \rightarrow K\bar{K}$ is dominated by the ϕ and should be close to the unitarity limit. On the other hand $K\bar{K} \rightarrow \pi \rho$ gets contributions from mainly ω resonances (or reggeons dual to these). Experimentally the ω 's lie far from the ϕ in units of their widths. Therefore the ω amplitude is small at the ϕ mass and we expect the ratio (3) to be small.[†]

[†] From a two channel resonance formalism one has

$$\left| \frac{A(K\bar{K} \rightarrow \rho\pi)}{A(K\bar{K} \rightarrow K\bar{K})} \right| = \frac{m_\omega (\Gamma_{\omega K\bar{K}} \Gamma_{\omega \rho\pi})^{\frac{1}{2}}}{m_\phi^2 - m_\omega^2},$$

where the widths are to be evaluated at the ϕ mass. Using an SU(3) value [8] $g_{\omega K\bar{K}}^2/4\pi = 1.11$ one gets $\Gamma_{\omega K\bar{K}} = 3$ MeV. From the experimental $\Gamma_{\phi \rho\pi} = 0.73$ MeV and $\phi \approx 0.06$ we have $\Gamma_{\omega \rho\pi} \approx 200$ MeV which agrees within a factor of 3 with predictions from $\omega \rightarrow \pi^0 \gamma$ and vector meson dominance as well as with the Gell-Mann-Sharp-Wagner model taking into account the finite ρ width. This gives a value for the above ratio of 0.04 which is to be compared with 0.06. Considering that higher ω resonances should increase the value as well as the approximate nature of eq. (4) the agreement is good.

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Thus we have consistency with an OZI suppression factor $\theta^\phi = 0.06$ discussed previously.

On the other hand if there were an ω resonance close to the ϕ in mass, the ratio in eq. (4) would have to be large and no OZI suppression could be consistent with unitarity. The ϕ and the hypothetical ω resonance would be strongly mixed.

For the psions similar arguments apply. Apart from phase space factors we get for the region $4m_D^2 < s < (m_D + m_{D^*})^2$:

$$\frac{\Gamma_{\text{n.h.}}^{\text{D}\bar{\text{D}}}}{\Gamma_{\text{D}\bar{\text{D}}}^{\text{D}\bar{\text{D}}}} \approx \frac{\sum_i |A(\psi \rightarrow i)|^2}{|A(\psi \rightarrow \text{D}\bar{\text{D}})|^2} \approx \frac{\sigma(\text{D}\bar{\text{D}} \rightarrow \text{n.h.})}{\sigma(\text{D}\bar{\text{D}} \rightarrow \text{D}\bar{\text{D}})} \quad (5)$$

Here the sum runs over all normal hadron channels. If the ω resonances in $\text{D}\bar{\text{D}} \rightarrow \text{n.h.}$ are not very much suppressed compared to the psi resonances in $\text{D}\bar{\text{D}} \rightarrow \text{D}\bar{\text{D}}$ the ratio in (5) is not very small. Alternatively one can use semilocal duality and approximate (5) by:

$$\frac{\sigma(\text{D}\bar{\text{D}} \rightarrow \text{n.h.})}{\sigma(\text{D}\bar{\text{D}} \rightarrow \text{D}\bar{\text{D}})} \propto \frac{\left(\frac{s - s_1}{s_0}\right)^{\alpha_\psi}}{\left(\frac{s - s_1}{s_0}\right)^{\alpha_\psi}} \quad (6)$$

which for large s falls rapidly assuming $\alpha_\psi < -3$. However, just above the $\text{D}\bar{\text{D}}$ threshold the scale factors s_0, s_1 are important.

With $s_1 \approx 4m_D^2$ one could not motivate a very small value for the inelasticity ratio in eq. (6). A third, more experimental, argument is to use $p\bar{p}$ annihilation near threshold as a guide. There the corresponding ratio between annihilation and elastic scattering is large.

Therefore we do not find it likely that the ratio in (5) could be as small as 0.01 or less as required if the psions above 3.9 GeV are to have a suppression parameter θ comparable with or smaller than that of the ψ' (c.f. eq. (3)). A perhaps reasonable guess would be that the ratio is comparable to that of the ϕ , or $\theta \approx 0.05 - 0.2$, which also lies above the first new threshold.

Such a comparatively weak suppression for the psions above 3.9 GeV may be welcome experimentally, because then the branching ratio to normal hadrons could be much larger than usually expected. If this branching ratio would be of the order 50% one could understand why the detection of charmed mesons has been so difficult and why the effects in e.g. the energy dependence of the K/π ratio are so weak.

Above the $K\bar{K}^*$ or $\text{D}\bar{\text{D}}^*$ threshold another dynamical mechanism is possible to provide the self-consistency. There the $K\bar{K}^*$ (or $\text{D}\bar{\text{D}}^*$) contribution can at least in part cancel the $K\bar{K}$ (or $\text{D}\bar{\text{D}}$) contribution. Dual models predict in fact such a cancellation, because the thresholds contribute with alternating signs. The twisting of a line involves a phase factor $e^{i\pi\alpha}$. A loop with two physical particles of spin J_1 and J_2 then has a sign:

$$S = (-1)^{J_1 + J_2} \quad (7)$$

relative to a planar amplitude. Thus $K\bar{K}$ and $K\bar{K}^*$ have opposite signs. Using duality diagrams one can derive [9] a similar factor ($S = C_1 C_2$) involving charge conjugation, which for most situations (except the A_1 trajectory) is equivalent to eq. (7). This cancellation mechanism has been used to explain the rapid energy dependence of $\pi^- p \rightarrow \phi n$. [10]

The f' meson lies above the $K\bar{K}^*$ threshold. It can thus be suppressed in two ways:

- i) A small ratio corresponding to that of eq. (4), i.e. small amplitudes $K\bar{K}, K\bar{K}^* \rightarrow 2n\pi$ compared to the elastic amplitudes at $\sqrt{s} = m_{f'}$,
- ii) Cancellations between $K\bar{K}$ and $K\bar{K}^*$ contributions.

The magnitude of the latter depends apart from phase space and form factors on the coupling constants involved

$$(g_{f'K\bar{K}} g_{f'K\bar{K}} \text{ vs. } g_{f'K\bar{K}^*} g_{f'K\bar{K}^*}).$$

Earlier we argued that because of unitarity and because the $D\bar{D}$ discontinuity very likely is large the psions just above the $D\bar{D}$ threshold should have a large branching ratio into normal hadrons. Then it is natural to ask why the psions below the threshold can be so narrow. Of course, below this threshold the $D\bar{D}$ discontinuity ($\propto \text{Im}A$) vanishes. A priori one expects the real part to be of the same order of magnitude as the imaginary part above the threshold. At the thresholds the real part should be continuous. Therefore it need not be very small below the threshold.

We present below one argument for why the real part is very small for the vector mesons. It is more model dependent than our other

results. Qualitatively it is related to the same cancellation mechanism as discussed previously: The real part is small because the dispersive contributions from intermediate states of positive sign (in particular $D\bar{D}$ and $D^*\bar{D}^*$) are cancelled by the contributions of negative signs ($D\bar{D}^* + D^*\bar{D}$).

We present a more comprehensive discussion elsewhere [11] and here only sketch the main arguments and results. First, we factorize from the amplitude of fig. 2a, factors containing the ψ and ψ' poles[†]:

$$A(s) = (1 + \frac{c}{s_{\psi}-s}) (1 + \frac{c'}{s_{\psi'}-s}) \bar{A}(s), \quad (8)$$

[†] This can be compared with

$$B_4(-\alpha_s, -\alpha_t) = (1 + \frac{\alpha_t}{\alpha_s}) (1 + \frac{\alpha_t}{1-\alpha_s}) B_4(2 - \alpha_s, -\alpha_t)$$

Assuming c, c' to be small the energy dependence of the psi residues will be given by $\bar{A}$.

The imaginary part of $\bar{A}$ gets contributions from all open channels $D\bar{D}, D\bar{D}^* + D^*\bar{D}, D^*\bar{D}^*$ etc.

$$\text{Im } \bar{A}(s) = \sum_{i=1}^N \text{Im } \bar{A}_i(s). \quad (9)$$

One can assume each $\text{Im } \bar{A}_1(s)$ to have a high energy behaviour given by a Regge cut involving a low lying D^* trajectory. Then each $\text{Im } \bar{A}_1(s)$ should after an initial rise at the threshold fall rapidly with a large negative power in s . Assuming the amplitude to be well behaved at infinity and that other singularities can be neglected we have a dispersion relation

$$\text{Re } \bar{A}(s) = \frac{1}{\pi} \int_{4m_D^2}^{\infty} \frac{\text{Im } \bar{A}(s')}{s' - s} ds' \quad (10)$$

which then gives a reasonable estimate of the real part in a narrow region below the threshold including the $\psi'(3.7)$ and perhaps the $\psi(3.1)$.

In fig. 3a we show a schematic picture of the situation assuming the masses to be: $m_D, m_{D^*}, m_{D^{**}} = 1.9, 2.1, 2.4$ GeV. The signs of $\text{Im } \bar{A}_1$ come from eq. (7). It is easy to convince oneself numerically that if the area of the $D\bar{D}^* + D^*\bar{D}$ peak is about twice in magnitude of that of the $D\bar{D}$ peak the contributions to the real part from these will cancel each other in the neighborhood of the ψ' . Other contributions from more distant singularities of course shift the position of the zero thus introduced. However, it is reasonable to assume that the true amplitude has a zero (or at least a minimum) near the ψ' . Then one can understand the smallness of $\Gamma^{n.h.}$ of ψ' and ψ , as well as why $\Gamma^{n.h.}(\psi') < \Gamma^{n.h.}(\psi)$.

Without this alternation in sign a similar argument shows that the $D\bar{D}$ and $D\bar{D}^*$ would together give a real part at the ψ' mass only a factor 2-3 smaller than the imaginary part above the threshold. With the cancellation one can understand an additional suppression factor of ≈ 10 in amplitude.

In fact, looking at the other spin states (fig. 3b and c) the cancellation works differently. Parity forbids the $D\bar{D}$ channel for the pseudoscalar and axial vector states and $D\bar{D}^*$ for the scalar mesons. Therefore we expect a much larger real part at least for the scalars because all nearest thresholds contribute with the same sign. Thus the scalars and possibly the pseudoscalars should be less pure $c\bar{c}$ states in agreement with most current expectations.

Another interesting decay is $\psi' \rightarrow \psi \pi \pi$. Notice that here the twists occur in a different way (fig. 2d). Therefore the signs of the contributions from different intermediate states in the loop are different. For example, let us assume the $\pi\pi$ state to be in a 0^{++} state (ϵ) and that the $D_1 \bar{D}_2 \rightarrow \psi \epsilon$ amplitude can be approximated by D and D^* particle exchanges. Then ordering the D states as in fig. 2d parity forbids $D_2 \bar{D}_3 = PV$ or VD which "could have cancelled" the allowed contributions $D_2 \bar{D}_3 = PP$ or VV which have the same sign. The sign does not depend on D_1 ; but for the nearest thresholds $D_1 = P$ or V . The "nearest" allowed contributions of opposite sign involve tensor mesons ($D_2 \bar{D}_3 = VT$ or TV). Of course this picture is modified when the exchanges in $D\bar{D} \rightarrow \psi \epsilon$ are reggeized and one allows for e.g. axial vector exchanges to cancel the PP contribution, but

still we find it unlikely that the cancellation mechanism in the real part is effective. Therefore we expect $\psi' \rightarrow \psi\pi\pi$ not to be suppressed by the cancellation and one can understand the large experimental branching ratio for $\psi' \rightarrow \psi\pi\pi$. The psions above the $D\bar{D}$ threshold should then have widths to $\psi\pi\pi$ only a few times that of ψ' (or ≈ 0.5 MeV), because without the additional suppression due to the cancellation in the real part the energy dependence should be much smoother than in fig. 1.

Sufficiently far below the $D\bar{D}$ threshold one can hardly neglect the singularities due to normal hadrons. It is reasonable to assume that because of these the OZI breaking again increases towards smaller s (c.f. the dashed part of the curve in fig. 1).

All considerations above apply of course equally well with D, D^* , etc. replaced by the $c\bar{s}$ states F, F^* , etc. These are important for $SU(3)$ and $SU(4)$ considerations. We expect that near the new thresholds $SU(3)$ is badly broken by a small mass difference $m_F - m_D$ because of the steep energy dependence of the OZI breaking.

The crucial experimental test to our arguments is the prediction that the psions above the $D\bar{D}$ threshold have widths to normal hadrons much larger than the $\psi(3.7)$, corresponding to branching ratios $x^{n.h.}$ larger than about 1%. In our picture $x^{n.h.}$ could be even $\geq 50\%$. In such a case one should at SPEAR see the psions above 3.9 GeV also in $3\pi, 5\pi$, etc. mass distributions. The rise in R at 3.9 GeV would then to a much smaller fraction than is usually believed be due to actually produced charmed mesons.

Acknowledgments

We thank Professors G. F. Chew, J. D. Jackson and Dr. P. Hoyer for useful comments and Professor J. D. Jackson for his kind hospitality at the Lawrence Berkeley Laboratory.

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FIGURE CAPTIONS

Fig. 1. A schematic picture of the energy dependence of the OZI breaking as measured by the parameter θ (eq. (1)). The data points are obtained from experimental ψ widths.

Fig. 2.(a),(b),(c). The lowest order OZI breaking diagram drawn in ways which are approximately equivalent by semilocal duality.

(d) The quark line diagram for $\psi' \rightarrow \psi \pi \pi$.

Fig. 3.(a) The contributions to the imaginary part due to $D\bar{D}$, $\bar{D}D^* + D\bar{D}^*$, $D^* \bar{D}^*$ etc. thresholds for the vector (tensor) $c\bar{c}$ states. The dashed line shows the behaviour of the real part below threshold.

(b) and (c). As in (a) but for the $O^-(1^+)$ and O^+ mesons respectively. Some of the thresholds are forbidden by parity.

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Fig. 1

Fig. 2

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(a) $\psi(3.7) \psi(3.95) \psi(4.1) \psi(4.4) \quad I^{--}(2^{++})$

(b) $O^{-+}(1^{+})$

(c) O^{++}

Fig. 3

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