THE OZI RULE VIOLATION IN $N\bar{N}$ ANNIHILATION AT REST AND THE SKYRME MODEL

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The question is raised if a large violation of the OZI rule, recently observed in many channels in $N\bar{N}$ annihilation at rest, can be explained in the framework of the Skyrme model.

1 Introduction

According to the OZI-rule, $\phi$-meson production in $N\bar{N}$ annihilation can happen only due to its nonstrange quark content and should be significantly suppressed:

$$R(X) = \frac{BR(N\bar{N} \rightarrow \phi X)}{BR(N\bar{N} \rightarrow \omega X)} \sim tg^2\delta = (2.5 \pm 1.5) \cdot 10^{-3} ,$$

where $X$ stands for the accompanying particle(s) and $\delta \approx 4^\circ$ is a small deviation from the ideal $\omega - \phi$ mixing angle.

This prediction is not confirmed by recent high-statistics experiments at LEAR. A large excesses over the naive OZI expectations were reported in the $\phi$-meson production in $N\bar{N}$ annihilation at rest by ASTERIX, Crystal Barrel and OBELIX collaborations. In the most distinguished channels the OZI rule violation reaches about two orders of magnitude:

$$R(\pi_0) = 0.14 \pm 0.04 , \quad R(\pi^-) = 0.16 \pm 0.04 , \quad R(\gamma) = 0.33 \pm 0.15 .$$

The weaker violation of the quark-line rule is also observed in many other channels.

The $\phi$-meson production in nucleon antinucleon annihilation at rest reveals the following features:

- There are strong enhancement over the naive OZI prediction for the channels $\phi\pi$, $\phi\gamma$. In the channels $\phi\rho$, $\phi\omega$, $\phi\pi\pi$ there are by one order of magnitude smaller enhancement.
- There is no enhancement at all in the $\phi\eta$ channel.
For nucleon antinucleon annihilation at rest the OZI rule violation is much more stronger than for $\pi p$ or $pp$ scattering and higher energy $p\bar{p}$ annihilation.

The large enhancement of $\phi\pi$ appears to be due to S-wave annihilation, P-wave annihilation exhibiting no large deviation from the OZI prediction.

Several models were suggested to explain the large $\phi$ production rate in $N\bar{N}$ annihilation at rest, neither of them being completely successful in incorporating the above mentioned features of the annihilation.

The most popular explanation of the OZI-rule violation assumes the existence of the strange sea quarks admixture in proton. There are indications also from other areas of the elementary particle physics that the nucleon wave function contains some amount of $s\bar{s}$ pairs already in the non-perturbative regime at large distances. In this picture, a large OZI-rule violation is interpreted in terms of the "shakeout" and "rearrangement" of an $s\bar{s}$ component of the nucleon wave function. Why different channels are so drastically unequally effected by this strange sea, remains a mystery.

Another model involves an exotic four-quark resonance, presumably dubious $C(1480)$ meson, to explain an enhancement in the S-wave $\phi\pi$ production and lack of enhancement in the P-wave channel. However even more stronger OZI-rule violation in the $\phi\gamma$ channel might not be due to this mechanism, because $\phi\gamma$ has positive $C$-parity in contrast to $\phi\pi^0$ and can not be produced via intermediate crypto-exotic $1^{--}$ resonance.

One more possibility to overcome the quark line rule is to assume that $\phi$ meson is produced due to the final state interaction of kaons formed in the OZI allowed process $p\bar{p} \rightarrow K^*\bar{K} \rightarrow K\bar{K}\pi \rightarrow \phi\pi$. But this rescattering model can not explain the absence of OZI rule violation in the $\pi\pi\phi$ channel despite the fact that $K^*\bar{K}^*$ final state is even more copious than $K^*\bar{K}$ in $N\bar{N}$ annihilation. The concrete calculations in this model also don’t reproduce experimental data very well, giving more moderate OZI rule violation than observed.

Theoretically the quark line rule is justified by the $\frac{1}{N_c}$ expansion which shows that non-valence $q\bar{q}$ pairs are suppressed in mesons. But for baryons, owing the fact that they contain $N_c$ valence quarks, non-valence $q\bar{q}$ pairs is expected to be $O(N_c)$ more important than in mesons. So for baryons the OZI rule has no theoretical background even in large $N_c$ limit.

On the other hand, the structure of QCD simplifies in the large $N_c$ limit, where it is equivalent to an effective theory of weakly interacting mesons, as was shown by Witten. We expect that this simplified effective theory also
reveals the OZI rule violation for baryons and thus is a good starting point to study the phenomenon.

2 The Skyrme model

In the low energy limit, the (approximate) $SU(3) \times SU(3)$ chiral symmetry of QCD is spontaneously broken and the $\gamma_5$-phase of quark fields have nonzero vacuum expectation value. Let us rewrite the QCD lagrangian separating the vacuum chiral phase, which can be considered as an external field

$$L = \bar{\psi} \exp \left\{-\frac{i}{2f_\pi} \gamma_5 \phi \lambda^a \right\} i \partial \gamma_5 \phi \lambda^a \psi + \cdots \quad \text{(1)}$$

where $A_{R\mu} = W^+ \partial_\mu W, A_{L\mu} = W \partial_\mu W^+$ and $W = \exp \left(\frac{i}{f_\pi} \phi \lambda^a \right)$. In fact, pseudoscalar mesons are excitations of just this $\gamma_5$-phase above the vacuum. More precisely, because mesons are quark-antiquark bound states, they should be associated with the chiral phase of the quark bilocal $\bar{q}_L q_R$, that is with $U = \exp \left(\frac{i}{f_\pi} \phi \lambda^a \right) = W^2$.

Integrating the quark (and gluon) degrees of freedom from (1), we end with the effective theory of the $U(x)$ field:

$$L_{eff.} = \frac{f_\pi^2}{4} \text{Sp}(\partial_\mu U)(\partial^\mu U^+ ) + \cdots \quad \text{(2)}$$

We can expect that in the low energy limit the other terms in (2), containing more derivatives, are less significant when the first one.

But (2) is a purely meson theory. Where are baryons? The crucial idea of T. Skyrme, revived in the context of QCD by E. Witten, is that we can still describe baryons by (2), considering them as vortexes of the mesonic "liquid".

Let us establish the most striking features of the Skyrme model, that these vortexes can carry a nonzero baryon number and a half-integer spin.

Maybe a nonzero baryon number, localized at Skyrmion, will appear not so strange if we remember about the negative energy Dirac’s sea of quarks. When the emergence of the baryon number can be considered as a "vacuum polarization" effect in the strong meson field of Skyrmion.

Because of the chiral anomaly, from (1) we have

$$\partial_\mu \bar{\psi}_R \gamma^\mu \psi_R = \frac{-1}{32\pi^2} \text{Sp} F_{\mu\nu}^R \tilde{F}_{\mu\nu} = \frac{-1}{8\pi^2} \partial^\mu \epsilon_{\mu\nu\lambda\sigma} \text{Sp}(A_{R\lambda}^\nu \partial_\sigma A_{R\lambda}^\nu + \frac{2}{3} A_{R\lambda}^\nu A_{R\lambda}^\nu A_{R\lambda}^\nu )$$

$$\partial_\mu \bar{\psi}_L \gamma^\mu \psi_L = \frac{1}{32\pi^2} \text{Sp} F_{\mu\nu}^L \tilde{F}_{\mu\nu} = \frac{1}{8\pi^2} \partial^\mu \epsilon_{\mu\nu\lambda\sigma} \text{Sp}(A_{L\lambda}^\nu \partial_\sigma A_{L\lambda}^\nu + \frac{2}{3} A_{L\lambda}^\nu A_{L\lambda}^\nu A_{L\lambda}^\nu )$$
But $A_\mu^R = W^+ \partial_\mu W$ and $A_\mu^L = W \partial_\mu W^+$ lead to

$$\varepsilon_{\mu \nu \lambda \sigma} \partial^\lambda A_\sigma^R = -\varepsilon_{\mu \nu \lambda \sigma} A_\lambda^R A_\sigma^R \quad \varepsilon_{\mu \nu \lambda \sigma} \partial^\lambda A_\sigma^L = -\varepsilon_{\mu \nu \lambda \sigma} A_\lambda^L A_\sigma^L$$

and $Sp(A_\nu^R A_\lambda^L A_\sigma^L) = -Sp(A_\nu^L A_\lambda^R A_\sigma^R)$. Thus

$$\partial_\mu \tilde{\psi}_R \gamma^\mu \psi_R = -\partial_\mu \tilde{\psi}_L \gamma^\mu \psi_L = \frac{1}{24\pi^2} \partial^\mu \varepsilon_{\mu \nu \lambda \sigma} Sp(A_\nu^R A_\lambda^L A_\sigma^L) .$$

(3)

For the baryon current $B_\mu = N_c \frac{1}{N_c} \tilde{\psi}_R \gamma^\mu \psi_R = \tilde{\psi}_R \gamma^\mu \psi_R$ (the first $N_c$ factor is due to color), (3) gives

$$\partial_\mu B_\mu = \partial_\mu (\tilde{\psi}_R \gamma^\mu \psi_R + \tilde{\psi}_L \gamma^\mu \psi_L) = \frac{1}{12\pi^2} \partial^\mu \varepsilon_{\mu \nu \lambda \sigma} Sp(A_\nu^R A_\lambda^L A_\sigma^L) =$$

$$= \frac{1}{24\pi^2} \partial^\mu \varepsilon_{\mu \nu \lambda \sigma} Sp[(U^+ \partial^\nu U)(U^+ \partial^\lambda U)(U^+ \partial^\sigma U)] ,$$

(4)

where we have introduced the meson field $U(x) = W^2(x)$ and the validity of the last step follows after a little algebra.

Equation (4) suggests the following expression for the baryon number current in terms of the meson field:

$$B_\mu = \frac{1}{24\pi^2} \varepsilon_{\mu \nu \lambda \sigma} Sp[(U^+ \partial^\nu U)(U^+ \partial^\lambda U)(U^+ \partial^\sigma U)] .$$

(5)

Inserting here the Skyrme’s hedgehog ansatz

$$U(x) = \exp \left\{ i \frac{\vec{r} \cdot \vec{r}}{r} F(r) \right\} ,$$

(6)

with the boundary conditions

$$F(0) = -n\pi \quad F(\infty) = 0 ,$$

(7)

we get after some computation (which is better to perform by some computer program, for example be REDUCE):}

$$B = \frac{1}{24\pi^2} \varepsilon_{ijk} \int Sp[(U^+ \partial_i U)(U^+ \partial_j U)(U^+ \partial_k U)] d\vec{x} =$$

$$= \frac{1}{2\pi^2} \int \frac{\sin^2 F}{r^2} \frac{dF}{dr} d\vec{x} = \frac{2}{\pi} \int_{F(0)}^{F(\infty)} \sin^2 F dF = n .$$

So the hedgehog (6) with $n = 1$ can be considered as a vortex with unit baryon number, that is as a nucleon. But to make such an identification, it should be shown that this hedgehog is a fermion, having a half-integer spin.
Of course the half-integer spin is the most marvelous thing which can be constructed from the spin zero pions. However this appears to be a rather general phenomenon, not only a peculiarity of the Skyrmion. In fact even a system from two spinless particles can have a half-integer spin if one of them is such a queer object as a magnetic monopole.

The motion of a charge $e$ in the field of the magnetic monopole $g$ is described by the equation ($r^2 = \vec{x} \cdot \vec{x}$)

$$m\ddot{x}_i = \frac{eg}{4\pi r^3} \varepsilon_{ijk} x_j \dot{x}_k.$$  

From this we find easily

$$\frac{d}{dt}[\varepsilon_{ijk} x_j (m\dot{x}_k)] = - \frac{d}{dt} \left[ \frac{eg}{4\pi} \frac{x_i}{r} \right].$$

Which indicates that for this system the angular momentum is

$$J_i = \varepsilon_{ijk} x_j (m\dot{x}_k) + \frac{eg}{4\pi} \frac{x_i}{r}.$$  \hfill (8)

Therefore the system possesses a half-integer spin and is a fermion for the lowest non-zero value $\frac{eg}{4\pi} = \frac{1}{2}$, allowed by the Dirac’s quantization condition.

This strange conclusion can be established on even more firm ground by showing that (8) is the Noether current density corresponding to an infinitesimal rotation.

As it is well known to find a singularity free Lagrangian for the charge-monopole problem is not a trivial task. Nevertheless a very elegant solution was found by Balachandran. By introducing SU(2) matrix $s$, defined as $X = \sigma_i \hat{x}_i = sps^+$, instead of the angular variables $\hat{x}_i = x_i/r$, a non-singular Lagrangian for the charge-monopole system can be written down in terms of $s$:

$$L = \frac{1}{2} m\dot{x}_i \dot{x}_i + i \frac{eg}{4\pi} Sp \{ \sigma_3 s^+ \dot{s} \} = \frac{1}{2} mr^2 + \frac{1}{4} mr^2 Sp \{ \dot{X}^2 \} + i \frac{eg}{4\pi} Sp \{ \sigma_3 s^+ \dot{s} \}.$$  \hfill (9)

Under a rotation we have $X' = \exp \{ i \frac{\varepsilon_i}{2} \sigma_i \} X \exp \{ -i \frac{\varepsilon_i}{2} \sigma_i \}$, that is $s' = \exp \{ i \frac{\varepsilon_i}{2} \sigma_i \} s$ and for an infinitesimal rotation

$$\delta s = i \frac{\varepsilon_i}{2} \sigma_i s , \quad \delta s^+ = -i \frac{\varepsilon_i}{2} s^+ \sigma_i , \quad \delta X = i \frac{\varepsilon_i}{2} [\sigma_i, X].$$  \hfill (10)
Using (10) we can find (note that $Sp\{\dot{X}[\sigma_i, \dot{X}]\} = 0$)
\[
\delta L = \dot{\epsilon}_i \left[ -\frac{i}{4} m r^2 Sp\{\dot{X}[\sigma_i, X]\} - \frac{eg}{24\pi} Sp\{X \sigma_i\} \right] = -\dot{\epsilon}_i \left[ \epsilon_{ijk} x_j (m \dot{x}_k) + \frac{eg}{4\pi} \hat{\epsilon}_i \right],
\]
and so the corresponding Noether current $-\frac{\delta L}{\delta \epsilon_i}$ coincides indeed with (8).

But what has all this to do with Skyrmions? Among the high-derivative terms, depicted by dots in (2), there is one which is analogous to the charge-monopole interaction term. This so called Wess-Zumino term has its root in the chiral anomaly and its contribution to the action looks quite exotic
\[
S_{WZ} = \left( \frac{i N_c}{240\pi^2} \right) \int d^5 x \, \epsilon_{\alpha\beta\gamma\delta\sigma} Sp\{(U^+ \partial^\alpha U)(U^+ \partial^\beta U)(U^+ \partial^\gamma U)(U^+ \partial^\delta U)(U^+ \partial^\sigma U)\}.
\]
Like the charge-monopole system, it is not possible to write down the corresponding piece of the (global) Lagrangian because the integral in (11) is over a 5-dimensional disc whose boundary is 4-dimensional Minkowskian space-time. But (11) simplifies enormously for the zero modes of the Skyrmion which are a time dependent $SU(3)$ rotations $s(t)$ of the hedgehog $U(r)$:
\[
U(r, t) = s(t) U(r) s^+(t).
\]
The piece of (11) involving $s$ looks like
\[
S_{WZ} = \left( \frac{i N_c B}{2\sqrt{3}} \right) \int dt \, Sp(\lambda s s^+ \dot{s})
\]
and exactly resembles the charge-monopole interaction term $\sim Sp(\sigma_3 s s^+ \dot{s})$. So the Wess-Zumino term produces a "monopole in $SU(3)$ space" and can lead to a half-integer spin just like how this happens in the charge-monopole system.

Note that the baryon number current (5) is just the Noether current corresponding to the singlet vector symmetry $U \rightarrow \exp \left( i \frac{\epsilon}{N_c} \right) U \exp \left( -i \frac{\epsilon}{N_c} \right)$ having the only non-zero contribution from the Wess-Zumino term! This is not surprising because both baryon number and the Wess-Zumino term have their origin in the chiral anomaly, as was sketched above. More amusing is that in the two flavor case the Wess-Zumino term vanishes while the baryon number, like the cheshire cat’s smile, still survives.

Thus the description of baryons as a mesonic vortexes is not so crazy as looks at first sight. Moreover the idea proved his fruitfulness in various applications, and we expect that it will be useful in the OZI rule violation studies also.
3 Coherent state picture of the $N\bar{N}$ annihilation

But first of all we need a description of the $N\bar{N}$ annihilation at rest as a Skyrmion anti-Skyrmion annihilation. Numerical calculations of this process have shown that just after the Skyrmion and anti-Skyrmion touch classical pion wave emerges as a coherent burst and takes away energy and baryon number as quickly as causality permits. This observation led R. Amado and collaborators to suggest the following simplified version of $N\bar{N}$ annihilation at rest.

After a very fast annihilation a spherically symmetric "blob" of pionic matter of size $\sim 1 \text{ fm}$, baryon number zero and the total energy twice the nucleon rest mass is formed. The further evolution of the system and the branching rates of various channels are completely determined by the parameters of this "blob" for which we can apply some simple phenomenological parametrization.

For example,

$$U(x) = \exp \{i \frac{\vec{r} \cdot \vec{F}}{r} \} , \quad F(r, t = 0) = \hbar \frac{r}{r^2 + a^2} \exp (-r/a) \quad (12)$$

where $\hbar$ is fixed by demanding that the total energy equals twice the nucleon mass and $a$ is a range parameter, the only free parameter of the model, assuming that the Skyrme model parameters are determined from the static properties of the nucleon at their usual values.

Starting from the initial field configuration (12), we can use the classical dynamical equations of motion to propagate the pion field, and other fields coupled to it, far away from the annihilation region there they no longer interact. These free classical radiation fields should be appropriately quantized because we detect particles in the final state and not the classical fields.

As was mentioned above numerical studies indicate that the resulting lump of the pionic matter propagates after the annihilation as a coherent burst of pion radiation. R. Amado et al. suggested that this peculiarity of the annihilation will be correctly represented if we assume that the asymptotic quantum state in the radiation zone is in fact a coherent state.

A coherent state $|\alpha>\$ is defined as an eigenstate of the annihilation operator

$$a|\alpha> = \alpha|\alpha> . \quad (13)$$

Using an usual commutation relation $[a, a^+] = 1$ and representing $|\alpha> = \sum_{n=0}^{\infty} c_n (a^+)^n |0>\$, where $|0>$ is the vacuum state, we find from (13) the recurrent relation $(n + 1)c_{n+1} = \alpha c_n$, from which it follows that up to normalization

$$|\alpha> = \exp (\alpha a^+) |0> .$$
It is not difficult to find out the normalization also. Let

\[ <\alpha | \alpha > = <0| \exp (\alpha^* a) \exp (aa^*) |0 > = f (\alpha^*, \alpha). \]

It is clear that \( f(0, \alpha) = f(\alpha^*, 0) = 1. \) On the other hand

\[
\frac{\partial}{\partial \alpha} f(\alpha^*, \alpha) = <0| \exp (\alpha^* a) a^+ \exp (\alpha a^+) |0 > = <0| (\frac{\partial}{\partial \alpha} \exp (\alpha^* a)) \exp (\alpha a^+) |0 > = \alpha^* f (\alpha^*, \alpha),
\]

so \( f(\alpha^*, \alpha) = \exp (\alpha^* \alpha) \) and the normalized coherent state looks like

\[ |\alpha > = \exp (-\frac{\alpha^* \alpha}{2}) \exp (\alpha a^+) |0 >. \] (14)

Let now \( |f > \) be the coherent state corresponding to the positive energy part of a quantum (scalar) field

\[
\varphi(\vec{r}, t) = \int \frac{d\vec{k}}{\sqrt{(2\pi)^3 2\omega_k}} (a_k e^{i\vec{k} \cdot \vec{r} e^{-i\omega_k t}} + a_k^+ e^{-i\vec{k} \cdot \vec{r} e^{i\omega_k t}}).
\]

That is for all \( \vec{k} \) we have \( a_k |f > = f(\vec{k}) |f > \). When

\[
<f | \varphi(\vec{r}, t) |f > = \int \frac{d\vec{k}}{\sqrt{(2\pi)^3 2\omega_k}} (f(\vec{k}) e^{i\vec{k} \cdot \vec{r} e^{-i\omega_k t}} + f^*(\vec{k}) f(\vec{k}) e^{-i\vec{k} \cdot \vec{r} e^{i\omega_k t}})
\]

is the associated classical field. So the generalization of (14) will be

\[ |f > = \exp \left(-\frac{1}{2} \int d\vec{k} f^*(\vec{k}) f(\vec{k}) \right) |0 >, \] (15)

where \( f(\vec{k}) \) is the Fourier transform of the classical radiation field.

The coherent state (15) has no definite 4-momentum. For pions it is also necessary to have a coherent state with a definite isospin. How to handle these subtleties the reader can find in the original literature\cite{28,30,31}.

Thus the logical scheme of the R. Amado et al.’s approach to the \( N\bar{N} \) annihilation at rest looks like this\cite{27}:

Given the initial (pion) field configuration and using the classical dynamical equations the asymptotic \( \pi, \rho, \omega \ldots \) fields should be generated. From those fields we construct one corresponding coherent state with definite 4-momentum and isospin. To find the probability of some final state we have to calculate the overlap of this state to the coherent state, which gives us the...
respective transition amplitude. The probability when is calculated as usual by integrating the absolute square of this transition amplitude over all final particle momenta and summing over all possible values of intermediate $\pi$ and $\rho$ isospin.

The physical reason why the above given picture of the $N\bar{N}$ annihilation is reasonable lies in the large number of produced particles. For the large final pion number its field can be approximated as classical. On the other hand if the number of quanta is very large an elimination of one of them doesn’t make a big difference. So the corresponding quantum state is in a good approximation an eigenstate of the annihilation operator.

Besides coherent states provide the closest quantum analog to the classical dynamics in the following sense. If we consider the standard position and momentum operators

$$q = \sqrt{\frac{\hbar}{2\omega}}(a^\dagger + a), \quad p = i\sqrt{\frac{\hbar\omega}{2}}(a^\dagger - a),$$

when in a coherent state $|\alpha> \rangle$ we have the corresponding uncertainties

$$(\Delta q)^2 = \langle |q|^2 |\alpha> \rangle - \langle |q| |\alpha> \rangle^2 = \frac{\hbar}{2\omega} \{[1 + (\alpha + \alpha^*)^2] - (\alpha + \alpha^*)^2\} = \frac{\hbar}{2\omega},$$

$$(\Delta p)^2 = \langle |p|^2 |\alpha> \rangle - \langle |p| |\alpha> \rangle^2 = -\frac{\hbar\omega}{2} \{[(\alpha^* - \alpha)^2 - 1] - (\alpha^* - \alpha)^2\} = \frac{\hbar\omega}{2}.$$

So the celebrated uncertainty relation is saturated for the coherent states.

In reality, however, the average number of pions produced in the $N\bar{N}$ annihilation at rest is about five with a variance of one, and so it is not a very large. Nevertheless a remarkably good agreement between calculated and measured characteristics of the annihilation was found by R. Amado et al.\textsuperscript{27, 28, 30, 31}. Note that the omega and rho mesons were also successfully included into the model despite the fact that their average numbers $\sim 1$.

4 The OZI rule violation

It is easy to see that the above described R. Amado et al.'s model of $N\bar{N}$ annihilation naturally incorporates the OZI rule violation for the $\phi$-meson production.
In the case of the ideal $\omega - \phi$ mixing, $\phi$-meson field decouples from the Skyrmion and is not directly excited by rotating hedgehog, which is believed to represent a nucleon in the collective coordinate approach. On the contrary $K$-meson field appears in the action with a linear coupling to the time derivative of the rigidly rotating hedgehog, because of the peculiar properties of the Wess-Zumino term and becomes excited.

So a nucleon in the Skyrme model carries some amount of the $K - \bar{K}$ cloud around it – the analog of the supposed strange component of the nucleon wave function, mentioned in the introduction. This means that after the annihilation initial field configuration contains not only pion field but also kaon (and antikaon) field(s). According to the spirit of the R. Amado et al.’s approach we have to evolve this kaon field using the classical equations of motion in the radiation zone and include in the coherent state. When from this coherent state we can extract the annihilation branching ratios for the kaon containing channels. But more important for us now is the excitation of the $\phi$-meson field due to its dynamical coupling to the $K$-meson field. While kaon field develops from the annihilation point to the radiation zone, $\phi$-meson field inevitably emerges because of obvious $K\bar{K} - \phi$ coupling and the coherent state describing annihilation should contain $\phi$-meson also, despite the fact that it is absent in the initial lump of the Skyrmion matter!

This transformation of the initial strangeness into the final $\phi$-meson resembles the rescattering model discussed in the introduction. Thus the Skyrme model picture of the $N\bar{N}$ annihilation at rest not only gives a natural framework for the OZI rule violation in the $\phi$-meson production, but also incorporates the characteristic features of two, at first glance very different, main models suggested to explain this phenomenon!

It is amusing that the OZI rule violation in the $\phi$-meson production appears to be one more manifestation of the Wess-Zumino term, because this very term is responsible for the initial kaon excitation in the nucleon which afterwards transforms into the $\phi$-meson.

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