# QUARK MODELS IN $N\overline{N}$ ANNIHILATION AT LOW ENERGIES

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*(Received July 19, 1994)*

Models for internal annihilation of quarks as a mechanism for baryon-antibaryon ($B_1\overline{B}_2$) production at low energy $p\bar{p}$ scattering are discussed. The results in form of branching ratios for different $B_1\overline{B}_2$ channels are compared with existing data either at a fixed center of mass energy $\sqrt{s}$ or at the same excess energy $\varepsilon = \sqrt{s} - m_{B_1} - m_{B_2}$. We find reasonable agreement with experimental data provided the internal quark interaction is of vector type. These results are contrasted to $q\bar{q}$ annihilation proceeding via an effective instanton interaction or in the context of Nambu–Jona-Lasinio type of models.

PACS numbers: 24.85. +p

## 1. Introduction

In recent years much effort has been devoted to understand the underlying mechanism of exclusive $p\bar{p}$ annihilation processes at low energies [1–9]. Experiments of this type have been done or are under way at the Low Energy Antiproton Ring (LEAR) at CERN. They can provide us with some clues to problems like OZI-rule violation [10] or the existence of exotic mesons [11]. However, in our opinion the most intriguing question so far is if we can describe the $p\overline{p}$ annihilation by using already established properties of the quark model. In this spirit we examine $p\overline{p}$ annihilation into two baryons using an internal annihilation model of quarks [2–3]. We emphasize here mainly two aspects of this model:

*(i)* In [12] it has been pointed out that a more sensible comparison with theory can be made if the experimental ratio of two cross sections (in [12] it was $\sigma(\Sigma^0\overline{\Lambda})/\sigma(\Lambda\overline{\Lambda})$) is taken not at the same center of mass energy $\sqrt{s}$ but at same excess energy $\varepsilon = \sqrt{s} - m_{B_1} - m_{B_2}$. This procedure of comparison can compensate better for the final/initial state interaction. We have performed the same analysis with more data points which is unfortunately limited due to crudeness of the data.

*(ii)* Beside the SU(6) wave functions the other ingredient to be specified in a quark model motivated picture of $p\overline{p}$ annihilation is of course the quark interaction itself. A natural way out would be to assume QCD. However, the low energy aspect of the collision processes we are interested in can cast some doubts on the reliability of perturbative QCD in this region. As we will see later in the context of our model the answer to this problem must remain open since no absolute prediction for the cross section are made. Hence the issue of the interaction must be decided by experiment. This is, however, inherent for all kind of models, not only for quark models. We would then like to stress that all possible interactions should be confronted with experimental data. We have performed this exercise for different kinds of interactions including an effective instanton [13] and Nambu–Jona-Lasinio Lagrangians [14]. It is then interesting to observe that the experimental data follow rather the conventional ideas of gluon-like annihilation or internal meson exchange than the more exotic ones.

Let us now outline what are the main motivations for constructing a quark model in a scattering process of hadrons at low energies. Speaking about the quark model one should be aware that the synonym "quark model" stands, in principle, for two different theoretical concepts. At high energies, in $p\overline{p}$ or deep inelastic $ep$ collisions, one usually speaks about current quarks with masses $m_u = 4$ MeV, $m_d = 7$ MeV, $m_s = 150$ MeV [15], whereas the low energy quark model invokes $m_u \simeq m_d \simeq 300$ MeV, $m_s = 500$ MeV and is called the constituent model. Needless to say that both concepts have been successful so far. Note that the natural place to probe parton model (*i.e.* current quarks) ideas is a collision process whereas the constituent model has proved successful in predicting static properties of hadrons such as magnetic moments of hyperons [16]. It is therefore quite natural to test also the constituent model in hadron- hadron reactions. That there is a theoretical gap in understanding the link between the two quark pictures supports only the need for such a test.

We would like to emphasize here that the meaning of such a test should be specified accurately. The theoretical predictions to be used in a true test of the constituent model should involve only established ideas of this model like *e.g.* SU(6) wave functions for hadrons. Recall that the original quark model was indeed based on symmetry groups like SU(3) and later SU(6) (to be precise we do not impose SU(6) symmetry on the $S$-matrix). It is this symmetry aspect of the quark model and the idea of internal quark annihilation [2–3] which we would like to confront with experimental data of $p\bar{p}$ collisions. In this case we have a clear idea about the expected accuracy of the predictions which is about 30% due to SU(6) breaking. Since SU(6) predictions given in terms of branching ratios like the famous $F/D$ value or magnetic moment ratios are quite reliable we do not expect a stronger SU(6) breaking in collision processes when comparing $\sigma(p\bar{p} \to B_1\overline{B}_2)/\sigma(p\bar{p} \to B'_1\overline{B'}_2)$. A systematic 30% deviation from experimental results would be even a stronger support for SU(6). Any usage of spatial dependent ($x$-dependent) wave function can obscure such a test since

*(i)* without additional input the quark model cannot provide us with such wave functions,

*(ii)* spatial wave functions are necessarily model dependent as the quark bound states are not fully understood.

It is then hard to judge if we are testing the SU(6) symmetry or the ansatz for the wave function itself. It should be clarified that we are not of the opinion that $x$-dependent wave functions should not be used at all. If they turn out to be universal *i.e.* independent of the specific process under consideration then indeed a deeper insight in the hadron structure can be gained [17]. However, in the first place, if we wish to compare the data with quark model predictions a more or less model independent comparison would be highly welcome. The price to be paid for neglecting spatial wave function is the lack of predictive power in angle and energy dependence of the cross sections. We therefore take here a modest point of view. The internal annihilation model outlined below is not intended to describe the whole complexity of $p\bar{p}$ physics at low energies. Its aim is rather to investigate if such a basic idea like internal $q\bar{q}$ annihilation projected on SU(6) wave functions is indeed supported by the data. In other words we put forward the question if out of the many reaction mechanisms, possible in $p\bar{p}$, the quark annihilation model can contribute substantially to the cross section and if, at which center of mass energies. Note that the reality can be more complex and the amplitude of $p\bar{p} \to B_1\overline{B}_2$ may be a coherent sum of several (unknown) reaction mechanisms (to show this explicitly we will compare the quark model with a meson exchange-like model). Intuitively one would expect that the $S$-wave annihilation of a $q\bar{q}$ pair [2–3] should contribute sizably to the cross section not far from the threshold. Anticipating our results below this indeed turns out to be the case. At higher energies higher partial waves start contributing, a fact which our model cannot accommodate for (here the usage of $x$-dependent wave functions is indeed unavoidable).

Our paper is organized as follows: In Section 2 we discuss internal fusion models based on $^3S_1$ (vector) and occasionally on $^3S_0$ (pseudovector) annihilation [2–3]. We will give a more detailed survey of this model for this type of interactions since all other types of interactions can be then treated in similar manner. A side remark on the production of charmed baryons is made. The last could be of relevance for Super–Lear. We omit a discussion of the other frequently used quark model which uses the $^3P_0$ (scalar) vertex [5, 8, 18]. In Section 3 we will show that using the concept of the excess energy yields indeed to a better agreement between experiment and theory. In Section 4 we address the question if more exotic type of interactions can also reproduce the data. We also comment briefly on a $K^*$ exchange model. Section 5 is left for conclusions.

## 2. The internal annihilation model

Before entering the technical details of the internal annihilation model let us mention a quantitative argument which supports this model. It is known [2] that the reactions

$$p\bar{p} \to \Xi\overline{\Xi}, \quad \Sigma^-\overline{\Sigma^-}, \quad \Sigma^-\overline{\Sigma^{*-}}, \quad \Sigma^{*-}\overline{\Sigma^{*-}}\ldots \tag{1}$$

are suppressed as compared to

$$p\bar{p} \to \Lambda\overline{\Lambda}, \quad \Sigma^0\overline{\Lambda}, \quad \Sigma^0\overline{\Sigma^0}\ldots. \tag{2}$$

This can be easily understood in terms of quark subprocesses since in the first case a double annihilation/creation is required.

At hand of the internal meson dominance [3] we will now outline the details of the model. The matrix elements to be computed are generically of the form

$$\langle p(\sigma)\bar{p}(\bar{\sigma})|\mathcal{L}(x)\mathcal{L}(x')|B_1(\tau)\overline{B}_2(\bar{\tau})\rangle\,, \tag{3}$$

where $\sigma$, $\bar{\sigma}$, $\tau$, $\bar{\tau}$ denote the spin degrees of freedom.

We will be concerned with all possible currents of the form

$$\overline{\Psi}\Gamma_A\Psi, \qquad \Gamma_A = 1, \gamma_\mu, \sigma_{\mu\nu}, \gamma_5\,, \tag{4}$$

where the physical picture behind the vector current is the gluon exchange. In this case expression (3) takes the form

$$\langle p(\sigma)\bar{p}(\bar{\sigma})|(\overline{\Psi}\lambda_c\gamma_\mu\Psi)(\overline{\Psi}\lambda_c\gamma^\mu\Psi)|B_1(\tau)\overline{B}_2(\bar{\tau})\rangle\,. \tag{5}$$

However, from technical point of view it is rather immaterial which current is sandwiched between the baryon states. Moreover it will turn out that with the exception of $n\overline{n}$, $\Delta\overline{\Delta}$ production the gluon annihilation and internal vector meson exchange will yield the same results. Hence it makes sense to start with the more general expression from which special cases can be easily deduced.

More explicitly the transition matrix element $S_{\mathrm{fi}}$ for internal meson exchange reads

$$\begin{aligned} S_{\mathrm{fi}} &= (-ig)^2 \sum_M \int d^3x \int d^3x' \Delta_M(x-x') \\ &* \langle p(\sigma)\overline{p}(\overline{\sigma})|\overline{\Psi}_{\xi,A,a}(x)\Psi_{\xi',A',a'}(x)\overline{\Psi}_{\eta,B,b}(x')\Psi_{\eta',B',b'}B|B_1(\tau)\overline{B}_2(\overline{\tau})\rangle \\ &* \Gamma_{AA'}\Gamma_{BB'}\delta^{\mathrm{c}}_{\xi\xi'}\delta^{\mathrm{c}}_{\eta\eta'}\delta^{\mathrm{fl}}_{aa'}\delta^{\mathrm{fl}}_{bb'}\ \langle q_a\overline{q}_a|M\rangle\langle M|Q_b\overline{Q}_{b'}\rangle\,, \end{aligned} \tag{6}$$

where the sum runs over ($\rho$, $\omega$, $\Phi$) mesons for the vector case ($\Gamma = \gamma_\mu$) and ($\pi^0$, $\eta$, $\eta'$) mesons for the pseudoscalar exchange ($\Gamma = \gamma_5$). The labels c and fl at the delta functions refer to colour and flavour, respectively. Since, as mentioned before, we neglect the energy dependence of the matrix elements the propagator $\Delta$ in momentum space can be written as

$$\Delta_M = i\left(-m_M^2 + im_M\Gamma_M\right)^{-1} \tag{7}$$

which is connected to $\Delta_M(x-x')$ through a Fourier transformation

$$\Delta_M(x-x') = \frac{1}{(2\pi)^4}\int d^4q\ \mathrm{e}^{-iq(x-x')}\Delta_M = \delta^{(4)}(x-x')\Delta_M\,, \tag{8}$$

Note that the gluon exchange model is obtained from (6) by the obvious replacement $\delta^{\mathrm{c}} \to \lambda^{\mathrm{c}}$ and by dropping the $M$ dependent couplings $< q\overline{q}|M >$.

In (6) we have spelled out explicitly the index structure of the matrix element since the calculation is simplified after performing a charge conjugation transformation and rearranging the indices according to Fierz transformation of the Dirac indices

$$\Gamma_{AA'}\Gamma_{BB'} = \sum_{\alpha=1}^{16} c_\alpha \Gamma^\alpha_{AB'}\Gamma^\alpha_{BA'} \tag{9}$$

and similar ones in the flavour (colour) space

$$\begin{aligned} \delta_{\alpha\alpha'}\delta_{\beta\beta'} &= \frac{1}{2}\sum_{A=1}^{8}\lambda^A_{\alpha\beta'}\lambda^A_{\beta\alpha'} + \frac{1}{3}\delta_{\alpha\beta'}\delta_{\beta\alpha'}\,, \\ \sum_{B=1}^{8}\lambda^B_{\xi\xi'}\lambda^B_{\eta\eta'} &= -\frac{1}{3}\sum_{B=1}^{8}\lambda^B_{\xi\eta'}\lambda^B_{\eta\xi'} + \frac{16}{9}\delta_{\xi\eta'}\delta_{\eta\xi'}\,. \end{aligned} \tag{10}$$

Going to the non-relativistic limit the relevant matrix element to be computed takes the form

$$\begin{aligned} T_{fi} = & - \langle p(\sigma)|\overline{\Psi}_{\xi,A,a}\Psi_{\eta',B',b'}|B_1(\tau)\rangle\langle p(\overline{\sigma})|\Psi_{\eta,D,b}\overline{\Psi}_{\xi',C,a'}|B_2(\overline{\tau})\rangle \\ & * \left[\tfrac{1}{2}\lambda^{\rm fl}_{ab'}\lambda^{\rm fl}_{ba'} + \tfrac{1}{3}\delta^{\rm fl}_{ab'}\delta^{\rm fl}_{ba'}\right]\left[\tfrac{1}{2}\lambda^{\rm c}_{\xi\eta'}\lambda^{\rm c}_{\eta\xi'} + \tfrac{1}{3}\delta^{\rm c}_{\xi\eta'}\delta_{\eta\xi'}{}^{\rm c}\right] \\ & * \left[\begin{pmatrix}1\\3\end{pmatrix}\delta_{AB'}\delta_{CD} + \begin{pmatrix}-1\\+1\end{pmatrix}\vec{\sigma}_{AB'}\vec{\sigma}_{CD}\right] \\ & * \sum_M \langle q_a\overline{q}_{a'}|M\rangle\langle M|Q_b\overline{Q}_{b'}\rangle\Delta_M\,, \end{aligned} \tag{11}$$

where the first number in the brackets refers to pseudoscalar, the second to vector meson exchange.

Since the matrix element (11) has to be evaluated in the Fock space it is rather convenient to use creation operators $a^{i+}_{q\alpha}$ ($\alpha$ is the spin index and $i$ denotes the colour degree of freedom) to represent the SU(6) baryon wave functions as follows

$$\begin{aligned} |p\uparrow\rangle &= \tfrac{1}{\sqrt{18}}\varepsilon_{ijk}\left\{a^{i+}_{u\uparrow}a^{j+}_{d\downarrow}a^{k+}_{u\uparrow} - a^{i+}_{u\downarrow}a^{j+}_{d\uparrow}a^{k+}_{u\uparrow}\right\}|0\rangle \\ |\Lambda\uparrow\rangle &= \tfrac{1}{\sqrt{12}}\varepsilon_{ijk}\left\{a^{i+}_{u\uparrow}a^{j+}_{d\downarrow}a^{k+}_{s\uparrow} - a^{i+}_{u\downarrow}a^{j+}_{d\uparrow}a^{k+}_{s\uparrow}\right\}|0\rangle \\ |\Sigma^0\uparrow\rangle &= \tfrac{1}{\sqrt{36}}\varepsilon_{ijk}\left\{2a^{i+}_{d\uparrow}a^{j+}_{s\downarrow}a^{k+}_{u\uparrow} - a^{i+}_{d\downarrow}a^{j+}_{s\uparrow}a^{k+}_{u\uparrow} - a^{i+}_{u\downarrow}a^{j+}_{s\uparrow}a^{k+}_{d\uparrow}\right\}|0\rangle \end{aligned} \tag{12}$$

...

Obviously the $\langle q\overline{q}|M\rangle$ coupling measures the $q\overline{q}$- content of the meson $M$. It is therefore necessary to comment on $\omega - \Phi$-mixing which we write in terms of a small parameter $\varepsilon$

$$\begin{aligned} |\omega\rangle &= \tfrac{1}{\sqrt{2}}\left[|u\overline{u}\rangle + |d\overline{d}\rangle\right] - \varepsilon|s\overline{s}\rangle \\ |\Phi\rangle &= \tfrac{-\varepsilon}{\sqrt{2}}\left[|u\overline{u}\rangle + |d\overline{d}\rangle\right] - |s\overline{s}\rangle\,. \end{aligned} \tag{13}$$

The value $\varepsilon^2 = 4\cdot 10^{-3}$ corresponding to quadratic Gell-Mann–Okubo mass formula is used consistently throughout the paper. One can, of course, introduce along the same line parameters $a$ and $b$ describing the mixing in the $\eta - \eta'$ sector. It will however turn out that the experimental data are in disagreement with the pseudoscalar fusion model whatever the values of the mixing parameters $a$ and $b$.

The final result of the calculation can be given in terms of a spin triplet amplitude $A^{(3)}$ and spin singlet part $A^{(1)}$

$$\sigma(B_1\overline{B}_2) \propto \frac{1}{4}\left\{3\left[A^{(3)}(B_1\overline{B}_2)\right]^2 + \left[A^{(1)}(B_1\overline{B}_2)\right]^2\right\}\,. \tag{14}$$

The calculation of $A^{(3)}$ and $A^{(1)}$ with the help of (11), (12) and (13) is lengthy but straightforward. The results are summarized in table 1 for the internal meson dominance model. The amplitudes $A^{(3)}$ and $A^{(1)}$ for the gluon exchange coincide with the vector case in Table I.

TABLE I

The amplitudes $A^{(3)}$ and $A^{(1)}$ entering Eq. (14) for the pseudoscalar and vector case denoted by a subscript ps. and v. respectively.

| Channel | $A^{(3)}_{\text{ps.}}$ | $A^{(1)}_{\text{ps.}}$ | $A^{(3)}_{\text{v.}}$ | $A^{(1)}_{\text{v.}}$ |
|---|---|---|---|---|
| $\Lambda\overline{\Lambda}$ | $0$ | $12$ | $12$ | $0$ |
| $\Sigma^0\overline{\Lambda}$ | $\frac{-4}{\sqrt{3}}$ | $0$ | $\frac{-8}{\sqrt{3}}$ | $\frac{-12}{\sqrt{3}}$ |
| $\Sigma^0\overline{\Sigma^0}$ | $\frac{8}{9}$ | $\frac{4}{3}$ | $\frac{28}{9}$ | $\frac{8}{3}$ |
| $\Sigma^+\overline{\Sigma^+}$ | $\frac{16}{9}$ | $\frac{8}{3}$ | $\frac{56}{9}$ | $\frac{16}{3}$ |
| $\Sigma^{*+}\overline{\Sigma^{*+}}$ | $\frac{8\sqrt{10}}{9}$ | $\frac{16}{3\sqrt{2}}$ | $\frac{-8\sqrt{10}}{9}$ | $\frac{-16}{3\sqrt{2}}$ |
| $\Sigma^+\overline{\Sigma^{*+}}$ | $\frac{16}{3\sqrt{18}}$ | $0$ | $\frac{-16}{3\sqrt{18}}$ | $0$ |
| $\Lambda\overline{\Sigma^{*0}}$ | $\frac{8}{\sqrt{6}}$ | $0$ | $\frac{-8}{\sqrt{6}}$ | $0$ |
| $\Sigma^{*0}\overline{\Sigma^{*0}}$ | $\frac{4\sqrt{10}}{9}$ | $\frac{8}{3\sqrt{2}}$ | $\frac{-4\sqrt{10}}{9}$ | $\frac{-8}{3\sqrt{2}}$ |
| $\Sigma^0\overline{\Sigma^{*0}}$ | $\frac{8}{9\sqrt{2}}$ | $0$ | $\frac{-8}{9\sqrt{2}}$ | $0$ |
| $n\overline{n}$ | $\frac{-32}{9}$ | $\frac{168}{9}$ | $\frac{104}{9}$ | $\frac{-96}{9}$ |
| $\Delta^{++}\overline{\Delta^{++}}$ | $\frac{8\sqrt{10}}{3}$ | $\frac{16}{\sqrt{2}}$ | $\frac{-8\sqrt{10}}{3}$ | $\frac{-16}{\sqrt{2}}$ |
| $\Delta^0\overline{\Delta^0}$ | $\frac{8\sqrt{10}}{9}$ | $\frac{16}{3\sqrt{2}}$ | $\frac{-8\sqrt{10}}{9}$ | $\frac{-16}{3\sqrt{2}}$ |

Note that in case of internal meson dominance model Eq. (14) must be multiplied by

$$\zeta_{\text{vector}} = \frac{\varepsilon^2 \Delta_{\Phi} + \Delta_\omega - \Delta_\rho}{2}\,,$$
$$\zeta_{\text{pseudosc.}} = \frac{-\Delta_{\pi^0} + a^2 \Delta_\eta + b^2 \Delta_{\eta'}}{2} \qquad (15)$$

for $n\overline{n}$ and $\Delta\overline{\Delta}$ production and by

$$\zeta_{\text{vector}} = \frac{\varepsilon}{\sqrt{2}}\left(\Delta_{\Phi} - \Delta_\omega\right)\,,$$
$$\zeta_{\text{pseudosc.}} = \frac{ab}{\sqrt{2}}\left(\Delta_\eta - \Delta_{\eta'}\right) \qquad (16)$$

for the other channels. For $a$ and $b$ see the explanation above.

TABLE IIa

Comparison of the theoretical predictions of internal pseudoscalar ($\sigma^{th.}_{ps.}$) and vector meson exchange ($\sigma^{th.}_{v.}$) with experimental data taken from [19]. Since the input is the cross section for $\Lambda\overline{\Lambda}$ at $\sqrt{s} = 3.6$ GeV (see Table IIb) all data at different center of mass energies have been conveniently rescaled. In case we have found two different experimental values at one and the same center of mass energy we consistently sticked to the lower value (see [3] where all data points have been listed).

| Channel | 2.4 GeV $\sigma_{exp.}$ $\times 77/127$ | 2.7 GeV $\sigma_{exp.}$ $\times 77/113$ | 3.0 GeV $\sigma_{exp.}$ $\times 77/91$ | $\sigma^{th.}_{ps.}$ | $\sigma^{th.}_{v.}$ |
|---|---|---|---|---|---|
| $\Lambda\overline{\Lambda}$ | $77 \pm 5$ | $77 \pm 10$ | $77 \pm 12$ | 77 | 77 |
| $\Sigma^0\overline{\Lambda}$ + c.c. | $39 \pm 3$ | $50 \pm 9$ | $58 \pm 10$ | 17 | 40 |
| $\Sigma^+\overline{\Sigma^+}$ | | $21 \pm 6$ | $30 \pm 5$ | 9 | 26 |
| $\Sigma^0\overline{\Sigma^0}$ | $7 \pm 3$ | 10 | 15 | 2 | 6 |
| $\Sigma^{*0}\overline{\Lambda}$ + c.c. | | | | 34 | 11 |
| $\Sigma^{*+}\overline{\Sigma^+}$ + c.c. | | | | 5 | 2 |
| $\Sigma^{*+}\overline{\Sigma^{*+}}$ | | | | 20 | 7 |
| $\Sigma^0\overline{\Sigma^{*0}}$ + c.c. | | | | 1 | 0.5 |
| $\Sigma^{*0}\overline{\Sigma^{*0}}$ | | | | 5 | 2 |
| $n\overline{n}$ | $1770 \pm 24$ | $1642 \pm 27$ | $1709 \pm 25$ | | 1727 |
| $\Delta^{++}\overline{\Delta^{++}}$ | $818 \pm 91$ | $1247 \pm 123$ | | | 1146 |
| $\Delta^0\overline{\Delta^0}$ | | | | | 127 |

It is experimentally known that $\Lambda\overline{\Lambda}$ is produced in a spin triplet state. As a consequence we must reject the pseudoscalar meson exchange. This conclusion is also supported by comparing the data with the theoretical results of the pseudoscalar fusion model (see Tables II a–b).

On the other hand Tables IIa and IIb yield a reasonable agreement between the experiment and the vector type exchange models for $p_{\overline{p}}$ momentum below 2.7 GeV. Due to crude data such a statement must be tantalizing but is in agreement with the expectation that at higher energies the unavoidable contribution of higher partial waves spoils the reliability of the S-wave $q\overline{q}$ annihilation model. Restricting the data to $p_{\overline{p}} \leq 2.7$ GeV has also the advantage that in this region the quoted experimental values [19] do not contradict each other. Then, however, little can be said about the $Y^*\overline{Y}$ and $Y^*\overline{Y}^*$ ($Y^*$ stands for decouplet members) production. Going therefore to $p_{\overline{p}} = 3.6$ GeV and $p_{\overline{p}} = 3.7$ GeV we see that $\sigma(\Sigma^{*+}\overline{\Sigma}^{*-})$ is still in agreement with the prediction. This is not a surprise since in this channel we are close enough to the threshold ($p_{\text{thr.}}(\Sigma^*\overline{\Sigma}^*) = 3$ GeV) to ensure an $S$-wave annihilation. The cross section for the $\Sigma^{*+}\overline{\Sigma^{+}}$ channel disagree with the prediction. At present we do not have an explanation for this failure and certainly more data are required to make a definitive statement.

TABLE IIb

The same comparison as in Table IIa but at higher center of mass energies. With the exception of the data point labelled by a asterisk which has been taken from [29] all other experimental values are from [19].

| Channel | 3.25 GeV $\sigma_{\text{exp.}}$ $\times 77/87$ | 3.6 GeV $\sigma_{\text{exp.}}$ | 3.7 GeV $\sigma_{\text{exp.}}$ $\times 77/82$ | $\sigma^{\text{th.}}_{\text{ps.}}$ | $\sigma^{\text{th.}}_{\text{v.}}$ |
|---|---|---|---|---|---|
| $\Lambda\overline{\Lambda}$ | $77 \pm 12$ | $77 \pm 20$ | $77 \pm 8$ | 77 | 77 |
| $\Sigma^0\overline{\Lambda}$ + c.c. | $50 \pm 10$ | $67 \pm 19$ | $65 \pm 9$ | 17 | 40 |
| $\Sigma^+\overline{\Sigma^+}$ | $32 \pm 12$ | $30 \pm 8$ | $41 \pm 8$ | 9 | 26 |
| $\Sigma^0\overline{\Sigma^0}$ | | 22 | 24 | 2 | 6 |
| $\Sigma^{*0}\overline{\Lambda}$ + c.c. | | $24 \pm 5$ | | 34 | 11 |
| $\Sigma^{*+}\overline{\Sigma^+}$ + c.c. | | $23 \pm 4^*$ | $34 \pm 10$ | 5 | 2 |
| $\Sigma^{*+}\overline{\Sigma^{*+}}$ | | $6.5 \pm 1$ | $5 \pm 2$ | 20 | 7 |
| $\Sigma^0\overline{\Sigma^{*0}}$ + c.c. | | | | 1 | 0.5 |
| $\Sigma^{*0}\overline{\Sigma^{*0}}$ | | | | 5 | 2 |
| $n\overline{n}$ | | $2000 \pm 600$ | | | 1727 |
| $\Delta^{++}\overline{\Delta^{++}}$ | $2434 \pm 221$ | $2020 \pm 160$ | $2723 \pm 282$ | | 1146 |
| $\Delta^0\overline{\Delta^0}$ | | | | | 127 |

It is also worthwhile observing that while the gluon exchange misses to explain the data on $n\overline{n}$ and $\Delta\overline{\Delta}$ production by orders of magnitude the internal meson fusion yields a reasonable agreement again at $p_{\overline{p}} \leq 2.7$ GeV. Since the ratio

$$\frac{\sigma(n\overline{n})}{\sigma(\Lambda\overline{\Lambda})} = \frac{3\left(\frac{104}{9}\right)^2 + \left(\frac{96}{9}\right)}{6 \times 144\varepsilon^2} \left|\frac{\Delta_\omega - \Delta_\rho}{\Delta_\omega - \Delta_\Phi}\right|^2 \tag{17}$$

is quite sensitive to the mixing parameter $\varepsilon$ (see Eq. (13)) we conclude that the same mechanism responsible for $\omega - \Phi$ mixing can also explain the difference of $s\overline{s}$ production as compared to $u\overline{u}$ and $d\overline{d}$. Since we do not know the mixing mechanism we have to rely on experimental information. Note that with the exception of $n\overline{n}$ and $\Delta\overline{\Delta}$ production it is the vector exchange which leads to reasonable agreement with data.

Let us now make a remark which could turn out to be important for the physics of the next low energy $p\bar{p}$ collider (Super–Lear). It is possible to enlarge the flavour sector of the produced hadrons and to consider the production of charmed baryons like $\Lambda_c^+\overline{\Lambda_c^+}$. There has been a recent revival of constructing relativistic wave functions for hadrons containing a heavy quark [24]. The static wave functions are assumed to be the usual non-relativistic SU(6)-type wave functions. The authors of [24] write for example the wave function for $\Lambda_c^+$ as (omitting colour indices)

$$\begin{aligned}|\Lambda_c^+\uparrow\rangle = \frac{1}{\sqrt{12}}\Big[ & u_\uparrow d_\downarrow c_\uparrow - u_\downarrow d_\uparrow c_\uparrow - d_\uparrow u_\downarrow c_\uparrow + d_\downarrow u_\uparrow c_\uparrow \\ & + c_\uparrow u_\uparrow d_\downarrow - c_\uparrow u_\downarrow d_\uparrow - c_\uparrow d_\uparrow u_\downarrow + c_\uparrow d_\downarrow u_\uparrow \\ & + d_\downarrow c_\uparrow u_\uparrow - d_\uparrow c_\uparrow u_\downarrow - u_\downarrow c_\uparrow d_\uparrow + u_\uparrow c_\uparrow d_\downarrow \Big] .\end{aligned} \tag{18}$$

It is then obvious that our previous results on internal $s\bar{s}$ production through gluon exchange can be carried over to the charmed sector without further modifications. It would be interesting to see if the agreement with experiment for charmed baryon production is as good as it was the case for the usual hyperons. We expect that charmed baryons will be produced at Super–Lear.

Other applications of the internal annihilation model can be found in [23, 25, 26].

## 3. Comparison with data at the same excess energy

Let us now introduce the concept of excess energy which is defined by

$$\varepsilon = \sqrt{s} - m_{B_1} - m_{B_2} . \tag{19}$$

Taking the ratio $\sigma_{B_1\overline{B}_2}/\sigma_{\Lambda\overline{\Lambda}}$ at the same excess energy means that given a measured point $\sigma_{B_1\overline{B}_2}(\sqrt{s})$ we are looking for $\sigma_{\Lambda\overline{\Lambda}}(\sqrt{s'})$ such that

$$\sqrt{s'} - 2m_\Lambda = \sqrt{s} - m_{B_1} - m_{B_2} . \tag{20}$$

In [12] it has been noted that by such a comparison we can partly get rid of final/initial state interaction. Our motivation to perform the analysis is the very good agreement between the theoretical prediction of $\sigma(\overline{\Lambda}\Sigma^0)/\sigma(\Lambda\overline{\Lambda})|_{\text{th.}} = 0.26$ and the corresponding experimental value 0.29 at a excess energy of $\varepsilon = 14.7$ MeV [12]. The analysis below is, however, restricted due to the fact that it is not always possible to find data points at two different center of mass energies such that Eq. (20) holds.

To account for phase space corrections we write

$$\sigma_{\text{ph.}}(B_1\overline{B}_2) = \sigma(B_1\overline{B}_2)\cdot\chi_{\text{ph.}}(B_1\overline{B}_2)\,,$$

$$\chi_{\text{ph.}}(B_1\overline{B}_2) = \frac{1}{s}\frac{\left[\left(s-(m_{B_1}+m_{B_2})^2\right)\left(s-(m_{B_1}-m_{B_2})^2\right)\right]^{1/2}}{\sqrt{s\,(s-4m_p^2)}} \tag{21}$$

and introduce for the theoretical prediction the branching ratio

$$R^{\text{ph.}}_{B_1\overline{B}_2} = \frac{\sigma^{B_1\overline{B}_2}_{\text{ph.}}(\sqrt{s})}{\sigma^{\Lambda\overline{\Lambda}}_{\text{ph.}}(\sqrt{s'})}\,. \tag{22}$$

Then, for instance for $\Sigma^0\overline{\Lambda}$ + c.c. we get (taking for a better orientation only the central values of the measured cross sections)

$$\frac{\sigma_{\Sigma^0\overline{\Lambda}+\text{c.c.}}(p_{\overline{p}}=2.4\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.2\,\text{GeV})} = 0.51, \quad R^{\text{ph.}}_{\Sigma^0\overline{\Lambda}+\text{c.c.}} = 0.46\,,$$

$$\frac{\sigma_{\Sigma^0\overline{\Lambda}+\text{c.c.}}(p_{\overline{p}}=2.7\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.43\,\text{GeV})} = 0.58, \quad R^{\text{ph.}}_{\Sigma^0\overline{\Lambda}+\text{c.c.}} = 0.47\,,$$

$$\frac{\sigma_{\Sigma^0\overline{\Lambda}+\text{c.c.}}(p_{\overline{p}}=3.0\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.7\,\text{GeV})} = 0.61, \quad R^{\text{ph.}}_{\Sigma^0\overline{\Lambda}+\text{c.c.}} = 0.47\,,$$

$$\frac{\sigma_{\Sigma^0\overline{\Lambda}+\text{c.c.}}(p_{\overline{p}}=3.2\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=3.0\,\text{GeV})} = 0.62, \quad R^{\text{ph.}}_{\Sigma^0\overline{\Lambda}+\text{c.c.}} = 0.48\,,$$

$$\frac{\sigma_{\Sigma^0\overline{\Lambda}+\text{c.c.}}(p_{\overline{p}}=3.7\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=3.45\,\text{GeV})} = 0.62, \quad R^{\text{ph.}}_{\Sigma^0\overline{\Lambda}+\text{c.c.}} = 0.48\,. \tag{23}$$

Again the agreement is worse at higher energies.

Similar analysis for $\Sigma^+\overline{\Sigma^+}$ and $\Sigma^{+*}\overline{\Sigma^{+*}}$ yields

$$\frac{\sigma_{\Sigma^+\overline{\Sigma^+}}(p_{\overline{p}}=2.7\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.2\,\text{GeV})} = 0.25, \quad R^{\text{ph.}}_{\Sigma^+\overline{\Sigma^+}} = 0.28\,,$$

$$\frac{\sigma_{\Sigma^+\overline{\Sigma^+}}(p_{\overline{p}}=3.0\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.43\,\text{GeV})} = 0.28, \quad R^{\text{ph.}}_{\Sigma^+\overline{\Sigma^+}} = 0.28\,,$$

$$\frac{\sigma_{\Sigma^+\overline{\Sigma^+}}(p_{\overline{p}}=3.25\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=2.7\,\text{GeV})} = 0.32, \quad R^{\text{ph.}}_{\Sigma^+\overline{\Sigma^+}} = 0.28\,,$$

$$\frac{\sigma_{\Sigma^+\overline{\Sigma^+}}(p_{\overline{p}}=3.6\,\text{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}}=3.15\,\text{GeV})} = 0.28, \quad R^{\text{ph.}}_{\Sigma^+\overline{\Sigma^+}} = 0.29 \tag{24}$$

and

$$\frac{\sigma_{\Sigma^{+*}\overline{\Sigma^{+*}}}(p_{\overline{p}} = 3.6\,\mathrm{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}} = 1.96\,\mathrm{GeV})} = 0.066, \quad R^{\mathrm{ph.}}_{\Sigma^{+*}\overline{\Sigma^{+*}}} = 0.045$$

$$\frac{\sigma_{\Sigma^{+*}\overline{\Sigma^{+*}}}(p_{\overline{p}} = 3.7\,\mathrm{GeV})}{\sigma_{\Lambda\overline{\Lambda}}(p_{\overline{p}} = 2.06\,\mathrm{GeV})} = 0.04, \quad R^{\mathrm{ph.}}_{\Sigma^{+*}\overline{\Sigma^{+*}}} = 0.046\,. \tag{25}$$

We remark that the agreement between theory and experiment in the other channels is not as good as in Eqs (23–25). This is partly due to the fact that the strength of initial/final state interaction can differ from channel to channel and partly to the rather modest phase space modification of the cross section in Eq. (21). Recall also that it is not always possible to find data points at two different lab momenta such that Eq. (20) is approximately fulfilled.

Considering the success of the two methods of comparison between theory and experiment we conclude that the $^3S_1$ annihilation model of quarks is very likely to be one of the main reaction mechanisms which governs the $p\overline{p}$ annihilation at low energies (provided we are not too far from the threshold). A more quantitative judgement, especially for each individual process, remains to be seen when more data are available.

## 4. Other interaction Lagrangians

Next we address the question if a more exotic quark interaction can also describe the data. In the first place we will discuss instanton interaction and then comment briefly on the pseudoscalar part of Nambu–Jona-Lasinio model.

Instantons play a significant role in the vacuum structure of QCD [20]. It has been pointed out by Dorokhov and Kochelev that quark interaction through instantons as introduced by 't Hooft [20] may play an important role in explaining hadron spectroscopy. In [21] the authors suggest that a large part of spin-spin splittings in hadron multiplets are not due to one gluon exchange but instead due to effective instanton interaction of quarks. The calculations performed in the framework of the bag model have shown that for instance a large part, if not the whole, of the $\Delta - N$ splitting can indeed be explained by the instanton induced quark-quark interaction. The rest of the splitting can then be accommodated by the second order of effective quark-quark interaction or by one gluon exchange. In this case we do not need to introduce a large quark-gluon coupling what has always been a problem in the bag model.

It is rather natural to investigate whether the effective instanton interaction that may be responsible for octet-decouplet mass splitting can also play a role in annihilation and creation of strange quarks in processes like $p\overline{p} \to Y\overline{Y}, Y\overline{Y^*}, Y^*\overline{Y^*}$.

Within the QCD vacuum model as an instanton liquid [13] the effective lagrangian for the four quark interaction is expressed as

$$\mathcal{L}^{(4)}_{\text{inst}} = \tfrac{2}{3}\pi^2\rho_c^2|\varepsilon_{ij}|\left[\overline{q}_{iR}q_{iL}\overline{q}_{jR}q_{jL}\left[1+\tfrac{3}{32}\left(1+\tfrac{3}{4}\sigma_i^{\mu\nu}\sigma^j_{\mu\nu}\right)\lambda_i^a\lambda_j^a\right]+(R\leftrightarrow L)\right], \tag{26}$$

where

$$q_{R,L} = \tfrac{1}{2}(1\pm\gamma_5)\,q \tag{27}$$

and $\rho_c = 1.6\,\text{GeV}^{-1}$ being the effective size of an instanton in the QCD vacuum. The indices $i, j$ in Eq. (26) refer to quark flavours. Here the convention for the antisymmetric tensor is $\sigma_{\mu\nu} = \frac{i}{2}[\gamma_\mu\,,\,\gamma_\nu]$.

Using the definitions of $q_{R,L}$ we can write $\mathcal{L}^{(4)}_{\text{inst}}$ in a more convenient form as follows

$$\begin{aligned}\mathcal{L}^{(4)}_{\text{inst}} = \text{const}|\varepsilon_{ij}|\Big[&\overline{q}_iq_i\overline{q}_jq_j + \overline{q}_i\gamma_5q_i\overline{q}_j\gamma_5q_j + \tfrac{3}{32}\Big(\overline{q}_i\lambda^aq_i\overline{q}_j\lambda^aq_j \\ &+ \overline{q}_i\gamma_5\lambda^aq_i\overline{q}_j\gamma_5\lambda^aq_j\Big) + \tfrac{9}{64}\left(\overline{q}_i\sigma_{\mu\nu}\lambda^aq_i\overline{q}_j\sigma^{\mu\nu}\lambda^aq_j\right)\Big].\end{aligned} \tag{28}$$

In contrast to Eq. (3) the relevant matrix element to be computed is now proportional to

$$\langle p(\sigma)\overline{p}(\overline{\sigma})|\mathcal{L}^{(4)}_{\text{inst}}|Y_1(\tau)\overline{Y}_2(\overline{\tau})\rangle\,. \tag{29}$$

In order to calculate this matrix element we will perform, as before, Fierz transformation in colour, flavour and spin space and finally take the non-relativistic limit. Note that according to Eq. (10) the term proportional to the small coefficient 3/32 in (28) will be multiplied by 16/9. Hence this term is not negligible as one might suspect at the first inspection of Eq. (28). The non-relativistic limit of (28) gives for the first and third term vanishing contribution. The second and the fourth term is proportional to spin zero projection operator $\frac{1}{2}(\delta\delta - \vec{\sigma}\vec{\sigma})$ whereas the fifth term goes with the spin-1 projection operator $\frac{1}{2}(3\delta\delta + \vec{\sigma}\vec{\sigma})$. Suppressing all indices we can write symbolically the final answer as

$$\mathcal{L}^{(4)}_{\text{inst}} \;\to\; \text{const}|\varepsilon_{ij}|\overline{\Psi}\,\overline{\Psi}\left[(\delta\delta-\vec{\sigma}\vec{\sigma}) + (3\delta\delta+\vec{\sigma}\vec{\sigma})\right]\Psi\Psi\,. \tag{30}$$

Comparing this expression with Eq. (11) we see that the instanton induced interaction results into a coherent sum of vector and pseudoscalar contributions.

To evaluate the matrix element (29) we use again static SU(6) wave functions for baryons. Our results are summarized in Table III. The peculiarity of the discussed instanton model is that it yields zero amplitudes for the production of decouplet members. The second problem is the equality of amplitudes in the singlet and triplet spin states. As mentioned before this is not in agreement with the experimental results on $\Lambda\bar{\Lambda}$ production. We can therefore safely state that the instanton model will not play a significant role in $p\bar{p}$ annihilation.

TABLE III

Theoretical predictions of the instanton model compared with experiment at $\sqrt{s} = 3.6$ GeV.

| Channel | $A^{(3)}_{\text{inst.}}$ | $A^{(1)}_{\text{inst.}}$ | $\sigma^{\text{th.}}_{\text{inst.}}$ | $\sigma_{\text{exp.}}(\sqrt{s} = 3.6$ GeV) |
|---|---|---|---|---|
| $\Lambda\bar{\Lambda}$ | 12 | 12 | 77 | $77 \pm 20$ |
| $\Sigma^0\bar{\Lambda}$ + c.c. | $\frac{-12}{\sqrt{3}}$ | $\frac{-12}{\sqrt{3}}$ | 25.7 | $67 \pm 19$ |
| $\Sigma^0\overline{\Sigma^0}$ | 4 | 4 | 8.6 | 22 |
| $\Sigma^+\overline{\Sigma^+}$ | 8 | 8 | 34.2 | $30 \pm 8$ |
| $\Sigma^{*+}\overline{\Sigma^{*+}}$ | 0 | 0 | 0 | $6.5 \pm 1$ |
| $\Sigma^+\overline{\Sigma^{*+}}$ + c.c. | 0 | 0 | 0 | $23 \pm 4$ |
| $\Lambda\overline{\Sigma^{*0}}$ | 0 | 0 | 0 | $24 \pm 5$ |
| $\Sigma^{*0}\overline{\Sigma^{*0}}$ | 0 | 0 | 0 | |
| $\Sigma^0\overline{\Sigma^{*0}}$ + c.c. | 0 | 0 | 0 | |

Any other Lagrangian can be treated in analogous manner. Consider for instance the scalar and pseudoscalar part of the four quark interaction in the Nambu–Jona-Lasinio model. The relevant interaction Lagrangian is given by [14]

$$\mathcal{L}_{\text{N-JL}} = \frac{G}{2}\left[\bar{\Psi}\lambda_a\Psi\bar{\Psi}\lambda_a\Psi - \bar{\Psi}\lambda_a\gamma_5 q\bar{\Psi}\lambda_a\gamma_5\Psi\right] . \tag{31}$$

Repeating the same procedure we have done for the instanton interaction one easily finds that (31) leads to the same results as the pseudoscalar internal meson exchange. Like the later we have to reject (31) in the context of low energy $p\bar{p}$ annihilation (see also [22]).

In the discussion of the instanton and NJ-L-Lagrangians we have omitted the inclusion of six quark interaction which in both models is in principle possible. For instance in the instanton model one has

$$\begin{aligned}\mathcal{L}^{(6)}_{inst} = &\frac{4\pi^2\rho_c^2}{3\langle 0|\bar{q}q|0\rangle}|\varepsilon_{ijk}|\bar{q}_{iR}q_{iL}\bar{q}_{jR}q_{jL}\bar{q}_{kR}q_{kL} \\ &\times\left\{1 + \tfrac{3}{32}\left[\lambda_i^a\lambda_j^a\left(1 + 3\vec{\sigma}_i\vec{\sigma}_j\right) + \tfrac{3}{10}d^{abc}\lambda_i^a\lambda_j^b\lambda_k^c\left(1 - 3\vec{\sigma}_i\vec{\sigma}_j\right) + \ldots\right]\right. \\ &\left. - i\tfrac{9}{64}f^{abc}\lambda_i^a\lambda_j^b\lambda_k^c\sigma_i^{\mu\nu}\sigma_j^{\nu\alpha}\sigma_k^{\alpha\mu}\right\} + (R \leftrightarrow L),\end{aligned} \tag{32}$$

where the dots indicate two other permutations. In the context of internal annihilation model we have to consider matrix elements of the form

$$\langle p(\sigma)\overline{p}(\overline{\sigma})|\mathcal{L}^{(6)}_{\text{inst}}|Y_1(\tau)\overline{Y}_2(\overline{\tau})\rangle \tag{33}$$

as higher order processes as compared to (29) (see however [13]). If it were not the case the calculation of (29) and even (33) would be meaningless as processes involving multi quark interaction like $\mathcal{L}^{(8)}$ *etc.* could turn out to be important. Note also that the simple factorization of Eq. (11) does not apply for $\mathcal{L}^{(6)}$.

Let us now go outside the quark model and compare the data with a simple meson exchange model. This is slightly outside the main focus of our paper, but as stated in the introduction such a comparison is indeed necessary. The relative importance of the various $K$-mesons in the reactions $p\overline{p} \to B_1\overline{B}_2$ can be judged by using the $\mathrm{SU}(3)_f$ symmetry for the coupling constants. According to this the $K(494)$ exchange is suppressed since we have

$$\frac{g^2_{KN\Sigma^0}}{g^2_{KN\Lambda}} = \frac{1}{27}\,. \tag{34}$$

On the other hand for the $K^*(892)$ coupling one gets

$$\begin{aligned} g^2_{K^*N\Lambda} : g^2_{K^*N\Sigma^+} : g^2_{K^*N\Sigma^0} &= 3 : 2 : 1 \\ g^2_{K^*N\Sigma^{*+}} : g^2_{K^*N\Sigma^{*0}} &= 2\,. \end{aligned} \tag{35}$$

To connect the octet and decouplet coupling constants we introduce a parameter $\alpha$ by setting $g^2_{K^*N\Lambda}/g^2_{K^*N\Sigma^{*+}} = 3/\alpha^2$. This parameter can be fixed by, for instance, the experimentally measured cross section $\sigma(p\overline{p} \to \Sigma^{*+}\overline{\Sigma^+} + \text{c.c.})$ which gives $\alpha = 0.58$ at $\sqrt{s} = 3.6$ GeV. Upon this we obtain the following ratios for the cross sections

$$\begin{aligned} &\sigma(\Lambda\overline{\Lambda}) : \sigma(\Lambda\overline{\Sigma^0} + \text{c.c.}) : \sigma(\Sigma^+\overline{\Sigma^+}) : \\ &\sigma(\Sigma^0\overline{\Sigma^0}) : \sigma(\Sigma^{*0}\overline{\Lambda} + \text{c.c.}) : \sigma(\Sigma^{*+}\overline{\Sigma^+} + \text{c.c.}) : \\ &\sigma(\Sigma^{*+}\overline{\Sigma^{*+}}) : \sigma(\Sigma^0\overline{\Sigma^{*0}}) : \sigma(\Sigma^{*0}\overline{\Sigma^{*0}}) \\ &= 77 : 51.3 : 34.2 : 8.6 : 17.3 : 23 : 3.9 : 5.8 : 1\,. \end{aligned} \tag{36}$$

This has to be compared with the experimental values in Table IIb at the center of mass energy 3.6 GeV. At the cost of an additional parameter the agreement is again reasonably good. It is then hard to decide which mechanism, the internal annihilation model or the $K^*$-exchange, is more important in $p\bar{p}$. However, bearing in mind the picture of the quark line diagrams for the quark annihilation it is quite likely that there exist some kind of duality between these two models. It is also tempting to seek for a deeper origin of the parameter $\alpha \simeq 1/\sqrt{3}$. We have checked that neither a relativistic version of SU(6) invariant lagrangian [27] nor the corresponding one for $SU(6)_W$ [28] gives the right answer. This is probably connected to the well-known problem to construct a consistent relativistic SU(6) symmetry.

## 5. Conclusions

The main concern of our paper has been to give a detailed survey of two aspects of the internal annihilation model in $p\bar{p}$ reactions. By using more data points we confirmed the previous observation that a comparison between experiment and theory is more sensible if performed at the same excess energy. Secondly we have put forward the question which interaction lagrangian is most suitable for the internal annihilation of quarks. The result of this analysis is that vector-like exchange yields the best agreement with experiment while instanton or NJ-L interactions will certainly not play a crucial role in $p\bar{p}$ at low energies.

We thank R. Decker for useful discussions and comments. This work has been supported in part by BMFT Grant Nr.06KA757 and by the Heraeus Foundation. M.N. wishes to thank Alexander von Humbold foundation for financial support under the Feodor-Lynen Fellowship program.

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