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TESTING QUARK REARRANGEMENT IN $N\bar{N}$ ANNIHILATION

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ABSTRACT

We use the constituent quark model to calculate s-wave $N\bar{N}$ scattering via quark rearrangement into three-meson states. With a quark Hamiltonian containing only a two-body confining interaction, we obtain a qualitatively correct description of the recent LEAR measurements. We can assign more than half of the s-wave annihilation cross-section to quark rearrangement. The results are, however, sensitive to the matching of the long-range meson exchange potential with the quark mediated annihilation potential.

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With the operation of the high intensity, high resolution antiproton facility LEAR, new data on the low-energy antinucleon-nucleon ($\bar{N}N$) interaction have appeared¹⁾. These data include measurements of annihilation, charge-exchange and elastic cross-sections down to beam momenta of 180 MeV/c. The relative importance of the various partial waves has been estimated from the elastic differential cross-section by a Legendre analysis. This gives the s-wave contributions to the annihilation and elastic cross-sections²⁾.

On the theoretical side, the application of quark cluster techniques has produced interesting new results in hadron-hadron scattering³⁾⁻⁶⁾. Simple quark interactions fitted to hadron spectroscopy very naturally generate interactions between hadrons described as colour-singlet quark clusters. For central $\bar{N}N$ collisions, there are two basic mechanisms in the quark dynamics. The first one consists of quark and antiquark rearrangement into $q\bar{q}$ meson and appears as a natural extension of quark interchange due to the Pauli principle in the exotic nucleon-nucleon^{3),4)} and K^+N interactions⁵⁾. The second mechanism is quark-pair annihilation or creation, already necessary for describing the decay of hadron resonances or non-exotic meson-baryon scattering in the region of resonances.

Both processes lead to meson final states in agreement with the dominance of the annihilation cross-section at low energy. The relative importance of rearrangement versus pair creation has been often discussed^{7),8)} by analyzing the relative branching ratios measured in $p\bar{p}$ annihilation at rest, but so far, no clear and unique conclusion has emerged. In this paper, we want to test the extreme hypothesis of pure quark rearrangement, cf., Fig. 1, by calculating its contribution to the complex optical potential describing s-wave $\bar{N}N$ scattering. We think indeed that it is not yet possible to treat rigorously quark-pair annihilation and creation in the non-relativistic quark model. For attempts in this direction, see Refs. 9) and 10).

In previous works^{7),11),12)} on quark rearrangement, the diagram of Fig. 1 was used as a guide to derive the shape and energy dependence of the optical potential, but room was left for phenomenological parameters to be adjusted by fitting the data. In contrast, in the present work, we remain in closer contact with the initial Hamiltonian and emphasize a consistent derivation of the $\bar{N}N$ interaction in the constituent quark model. We therefore do not attempt a very detailed description of the experimental data and restrict ourselves to the gross features.

The starting point of our investigation is the constituent quark model Hamiltonian

$$\mathcal{H} = \sum_{i=1}^6 \frac{\vec{p}_i^2}{2m_q} + \sum_{i < j} V_{ij}^{conf} - T_{cm} \quad (1)$$

The momenta $\vec{p}_i$ and positions $\vec{x}_i$ are numbered from one to three for the quarks, and from four to six for the antiquarks. T_{cm} is the centre-of-mass kinetic energy. The confinement interaction is chosen of the usual type

$$V_{ij}^{conf} = -a \tilde{\lambda}_i \tilde{\lambda}_j (|\vec{x}_i - \vec{x}_j|^n - A) \quad (2)$$

where the colour operator $\tilde{\lambda}_i \tilde{\lambda}_j$ between quarks is to be replaced by $-\tilde{\lambda}_i^T \tilde{\lambda}_j$ for antiquark-quark interactions. We will consider both the $n = 2$ case and the more realistic $n = 1$ linear potential. The two-body, colour-octet exchange character of potential (2) is a standard choice in multi-quark calculations. For a discussion and further references in particular on the problem of Van-der-Waals forces, see, for instance, Refs. 6) and 13).

Note that we will not consider in the following the effects of the QCD-residual (chromomagnetic) interactions, which split meson and baryon masses. The main reason lies in the complexity of the resulting $\bar{N}N$ to three meson transition potential. We will therefore use for the whole calculation average baryon and meson masses: $m_B = (m_N + m_\Delta)/2 = 1085$ MeV and $m_M = (3m_\rho + m_\pi)/2 = 611$ MeV. The co-ordinates relevant for the $\bar{N}N$ and three-meson (3M) channels are

$$\begin{aligned} \vec{\xi}_1 &= \vec{x}_1 - \vec{x}_2 & \vec{\xi}_2 &= \frac{\vec{x}_1 + \vec{x}_2 - 2\vec{x}_3}{2} \\ \vec{\xi}_3 &= \vec{x}_4 - \vec{x}_5 & \vec{\xi}_4 &= \frac{\vec{x}_4 + \vec{x}_5 - 2\vec{x}_6}{2} \\ \vec{r} &= \frac{\vec{x}_1 + \vec{x}_2 + \vec{x}_3 - \vec{x}_4 - \vec{x}_5 - \vec{x}_6}{3} \end{aligned} \quad (3)$$

and

$$\begin{aligned}\vec{\eta}_i &= \vec{x}_i - \vec{x}_{i+3}, & \vec{e}_i &= \frac{\vec{x}_i + \vec{x}_{i+3}}{2} \quad (i=1,2,3) \\ \vec{x} &= \vec{e}_1 - \vec{e}_2, & \vec{y} &= \frac{\vec{e}_1 + \vec{e}_2 - 2\vec{e}_3}{2}\end{aligned}\quad (4)$$

whereas $\vec{g} = \Sigma \vec{x}_i / 6$ denotes the centre-of-mass. The spatial part of the meson wave function is written as

$$\psi_{Mi} = f(\vec{\eta}_i) = (2\pi b^2)^{-3/4} \exp(-\vec{\eta}_i^2 / 4b^2) \quad (5)$$

and that of the baryons

$$\begin{aligned}\phi_N &= f(\vec{\xi}_1) f(\sqrt{4/3} \cdot \vec{\xi}_2) = (2\pi b^2)^{-3/4} (3\pi b^2/2)^{-3/4} \\ &\quad \cdot \exp(-\vec{\xi}_1^2 / 4b^2 - \vec{\xi}_2^2 / 3b^2).\end{aligned}\quad (6)$$

$$\phi_{\bar{N}} = f(\vec{\xi}_3) \cdot f(\sqrt{4/3} \cdot \vec{\xi}_4).$$

For simplicity, we use the same oscillator parameter b for both mesons and baryons, because this gives simple shell model type expressions for the $N\bar{N}$ and three meson system in the individual quark co-ordinates. We choose a constituent quark mass $m_q = 330$ MeV and determine the oscillator parameter $b = 0.51$ fm (related to the strength a) and the subtraction constant A by reproducing the ground state average masses m_B and m_M .

The intercluster wave function is described in the generator co-ordinate representation in both $N\bar{N}$ and three meson (3M) channels. Namely, generator co-ordinates $\vec{R}$, $\vec{X}$ and $\vec{Y}$ are associated to the relative co-ordinates $\vec{r}$, $\vec{x}$ and $\vec{y}$ defined in Eqs. (3) and (4) and the trial wave function is written as the following superposition

$$\psi = \mathcal{A} \left\{ \int c(\vec{R}) \psi_{\vec{R}}^{N\bar{N}} d^3\vec{R} + \sum_i \int d_i(\vec{X}, \vec{Y}) \phi_{\vec{X}, \vec{Y}}^{3M} d^3\vec{X} d^3\vec{Y} \right\} \quad (7)$$

where the operator $\mathcal{A}$ antisymmetrizes separately the quarks and the antiquarks. The cluster wave functions have the form

$$\psi_{\vec{R}} = \phi_N \phi_{\bar{N}} f[\sqrt{3}(\vec{r}-\vec{R})] f(\sqrt{12}\cdot\vec{g}) \chi_{N\bar{N}}^{SIC} \quad (8)$$

and

$$\phi_{\vec{x}\vec{y}} = \psi_{M_1} \psi_{M_2} \psi_{M_3} \cdot f[\sqrt{2}(\vec{x}-\vec{X})] f[\sqrt{8/3}(\vec{y}-\vec{Y})] f[\sqrt{12}\cdot\vec{g}] \chi_{3M}^{SIC} \quad (9)$$

and contain a spin-isospin-colour (SIC) part, the hadron wave functions of Eqs. (5) and (6), an intercluster distribution described by the same universal Gaussian shape and a centre-of-mass part peaked around the origin.

From the variational form of the Schrödinger equation $\langle \delta\psi | H-E | \psi \rangle = 0$, we obtain two coupled equations for the weight functions $c(\vec{R})$ and $d(\vec{X}, \vec{Y})$

$$0 = \int d\tau d^3\vec{R} \Psi_{\vec{R}}^* (H-E) \Psi_{\vec{R}} c(\vec{R}) + \sum_i \int d\tau d^3X d^3Y \Psi_{\vec{R}}^* (H-E) \phi_{\vec{x}\vec{y}}^i d_i(\vec{x}, \vec{y}) \quad (10)$$

$$0 = \int d\tau d^3\vec{R} \phi_{\vec{x}, \vec{y}}^{i*} (H-E) \Psi_{\vec{R}} c(\vec{R}) + \sum_i \int d\tau d^3X d^3Y \phi_{\vec{x}, \vec{y}}^{j*} (H-E) A \phi_{\vec{x}\vec{y}}^i d_i(\vec{x}, \vec{y}) \quad (11)$$

In the $N\bar{N}$ channel, both quarks and antiquarks have the correct symmetry. In the $3M$ channel, this is more delicate. The antisymmetrizer A can be rewritten as the product of a quark antisymmetrizer and a meson symmetrizer, say, $A = A_q S_M$. S_M acts on the co-ordinates $\vec{x}$ and $\vec{y}$ and forces the weight function $d(\vec{x}, \vec{y})$ to have the proper permutation symmetry. A_q would generate genuine meson-meson interactions by quark exchange which we cannot handle in the following. The action of A_q is therefore neglected.

The variational derivation of Eqs. (10) and (11) involves the transition matrix elements of the full Hamiltonian, and in particular, the non-orthogonality of the $N\bar{N}$ cluster wave function $\psi_{\vec{R}}$ and the three meson wave function $\phi_{\vec{x}\vec{y}}$ plays an important role. It has been argued^{7), 11)} that $N\bar{N}$ states are orthogonal to three meson states because baryons carry some "junction line" which has been also described as string centre in dual models. In such an approach, one needs a model for estimating the transition from two Y-shape flux tubes to three single tubes, otherwise one cannot derive the absolute strength for $N\bar{N} \rightarrow$ mesons. In our

opinion, it is preferable to avoid first such an extra assumption and push as far as possible the simple version of the quark model used in hadron spectroscopy. In fact, we have no evidence that the Hamiltonian of Eq. (1) is correct for multi-quark systems. But we think that together with the hypothesis of colour zero clusters, it presents a good starting point before adding complexities due to string breaking and recombination. The elimination of the meson channels from Eqs. (10) and (11) is possible for resonating group wave functions, where $\chi_{\xi, \eta}^{\text{RGM}}(\vec{x}, \vec{y}) = \delta^3(\vec{x} - \vec{\xi}) \delta^3(\vec{y} - \vec{\eta})$ instead of Gaussian functions (GCM) employed up to now in $\phi_{\vec{X}, \vec{Y}}^{\vec{x}, \vec{y}}(\vec{x}, \vec{y})$. By an additional defolding we derive an effective non-local NN annihilation potential¹⁴⁾:

$$\mathcal{V}_A(\vec{r}, \vec{r}') = \int_{\text{Mesons}} \mathcal{V}(\vec{r}, \vec{x}, \vec{y}) \mathcal{P}_r(\vec{x}, \vec{y}, \vec{x}', \vec{y}') \mathcal{V}(\vec{r}', \vec{x}', \vec{y}') \quad (12)$$

with the three-meson propagator

$$\mathcal{P}_r(\vec{x}, \vec{y}, \vec{x}', \vec{y}') = 6 \cdot \int \frac{d^3 p_x d^3 p_y}{(2\pi)^6} \frac{e^{i\vec{p}_x(\vec{x} - \vec{x}')} e^{i\vec{p}_y(\vec{y} - \vec{y}')}}{\frac{p_x^2}{2m_q} + \frac{p_y^2}{8/3 m_q} + 3m_H - (E_{\text{rel}} + 2m_{\frac{1}{3}} + i\epsilon) - \frac{3i\Gamma_H}{2}} \quad (13)$$

and the coupling potential

$$\mathcal{V}(\vec{r}, \vec{x}, \vec{y}) = N(\vec{r}, \vec{x}, \vec{y}) \cdot \frac{1}{\sqrt{3}} \cdot \left(\frac{1}{4m_q b^2} \left[4 - \frac{2}{3b^2} \left(\frac{3}{2} \vec{r}^2 + \vec{x}^2 + \frac{4}{3} \vec{y}^2 \right) - \frac{1}{3} \frac{\vec{k}_{\text{rel}}^2}{3m_q} \right] \right) \quad (14)$$

The overlap kernel is defined as

$$N(\vec{r}, \vec{x}, \vec{y}) = \delta \left(\frac{1}{2\pi b^2} \right)^{9/4} \exp \left[- \left(\frac{3}{2} \vec{r}^2 + \vec{x}^2 + \frac{4}{3} \vec{y}^2 \right) / 2b^2 \right] \quad (15)$$

The result of Eq. (14) agrees with the corresponding results of the Helsinki group^{7), 11)} in the leading term $V = N \cdot (1/\sqrt{3} m b^2)$. The other terms come from the kinetic energies in the Hamiltonian $\Sigma p_i^2 / 2m - T_{\text{cm}}$, due to the non-orthogonality

of the $\bar{N}\bar{N}$ and $3M$ wave functions. We have left out the summation over the various meson channels in Eq. (14). Since our Hamiltonian of Eq. (1) has full SU4-invariance, the summation over the complete set of SU4 meson states gives unity, i.e., $\sum_{3M} |\langle \bar{N}\bar{N} | 3M \rangle|^2 = 1$ for each spin-isospin channel. We allow one phenomenological parameter in the meson propagator $P_{\Gamma}(\vec{x}, \vec{y}, \vec{x}', \vec{y}')$ by replacing the real meson masses m_M by complex $m_M - i\Gamma_M/2$ where Γ_M is typically the ρ -meson width $\Gamma_M \approx 150$ MeV. From the Clebsch-Gordan coefficients giving the probabilities for the various meson channels, it is obvious that the vector mesons dominate the three meson channels, justifying our approximation. The quadratic qq-interaction [$n = 2$ in Eq. (2)] together with the harmonic oscillator wave functions has produced the special coupling matrix of Eq. (16), which leads to a simple separable potential. A linear confinement potential ($n = 1$) would not change that result. If we add an hyperfine interaction like $\vec{\sigma}_1 \vec{\sigma}_2 \vec{\lambda}_1 \vec{\lambda}_2 \delta(\vec{r}_{ij})$ to the Hamiltonian, we would get more complicated non-factorizable forces of the coupling potential with cross-terms in $\vec{r}\vec{x}$ and $\vec{r}\vec{y}$. Still in this case, one can see that the resulting potential acts only in even $\bar{N}\bar{N}$ -partial waves. With the $\vec{r}_{ij}^2$ interaction alone, quark rearrangement only gives an s-wave $\bar{N}\bar{N}$ interaction. It can be shown that relativistic corrections of $O(p^2/m_q^2)$ to scalar confinement produce a p-wave $\bar{N}\bar{N}$ interaction by allowing transition to ground state $q\bar{q}$ mesons in a relative $\ell = 1$ wave function. There is also the possibility for p-wave $\bar{N}\bar{N}$ -interactions to annihilate into higher mass p-wave mesons in a relative s-wave. Evaluation of the Eqs. (12)-(15) yields the optical potential from quark rearrangement in the following form;

$$V(\vec{r}, \vec{r}') = -6 \cdot \left(3 / 2^{11} \cdot \pi^3 \right)^{1/2} (1/b)^3 (1/m_q b^2) \cdot e^{-\frac{(3/4b^2)(r^2+r'^2)}{2}} \cdot f_{(16)}$$

$$\begin{aligned} f(E, \Gamma, \vec{r}, \vec{r}') = & \alpha(r) \alpha(r') \cdot I(E, \Gamma) - \frac{2}{3} (\alpha(r) + \alpha(r')) (2 + i\gamma I(E, \Gamma)) \\ & + \frac{4}{9} (6 - 2t + 2i\gamma - \gamma^2 \cdot I(E, \Gamma)) \end{aligned} \quad (17)$$

where

$$I(E, \Gamma) = \int_0^\infty ds \frac{e^{-s} s^2}{s - t - i\gamma}$$

All energies have been rescaled by $2m_q b^2$, e.g., the available kinetic energy for decay into three mesons is $t = 2m_q b^2 T = 2m_q b^2 (E_{\text{rel}} + 2m_B - 3m_M)$ and the width $\gamma = 2m_q b^2 \Gamma$. The form factor of the separable potential (14) is given by the exponential overlaps [cf. Eq. (13)] and functions α which are

$$\alpha(r) = r^2/b^2 - \frac{2}{3} (2m_B - 3m_M) \cdot 2m_q b^2 \quad (18)$$

In the limit where $\Gamma \rightarrow 0$ and $\alpha(r) \approx 1$, one can elucidate the resulting imaginary potential from quark rearrangement by writing it as a product of factors of different physical origin:

$$\begin{aligned} \text{Im } V(\vec{r}, \vec{r}') \approx & 6 \cdot \frac{1}{(2\pi b^2)^{3/2}} e^{-\frac{(\vec{r}-\vec{r}')^2 \cdot (3/8b^2)}{}} \cdot e^{-\frac{(\vec{r}+\vec{r}')^2 \cdot (3/8b^2)}{}} \cdot \left(\frac{4\pi}{3} b^3\right)^2 \\ & \cdot e^{-2m_q b^2 T} \cdot (-i\pi) \frac{(2m_q)^3 T^2}{8\pi^3} \cdot \left| \sum \bar{V}_{qq} \right|^2 \end{aligned} \quad (19)$$

The factors in Eq. (19) can be interpreted in the following way:

The imaginary potential is proportional to the number of possible reorderings of three quarks and three antiquarks. Its non-locality and range are determined by the hadron size (b). The propagation of the mesons takes place over a volume of $(4\pi/3)b^3$ and there are two independent degrees of freedom in the three-body system. The matching of wave functions of quarks in the $\bar{N}N$ sector with the meson wave functions becomes exponentially small ($e^{-2m_q b^2 T}$) with increasing available kinetic energy of the mesons. Finally, the imaginary potential is proportional to the transition matrix element squared which is of the order of magnitude of $\langle \vec{\lambda}_i \vec{\lambda}_j a | \vec{r}_i - \vec{r}_j |^2 \rangle \approx (1/3m_q b^2)$, and is multiplied by the phase space of the three-meson states of masses $m_M \approx 2m_q$ at an available kinetic energy T . Indeed comparing Eq. (19) to Eq. (16), in the limit $\gamma \rightarrow 0$ and $\alpha(r) \approx \alpha(r') \approx 1$

$$\text{Im } V_{(\vec{r}, \vec{r}')}^{\text{limit}} = (-0.04 i \pi) (2m_q b^2 T)^2 e^{-2m_q b^2 T} \cdot \frac{1}{b^3} \cdot \frac{1}{m_q b^2} \cdot e^{-\frac{3}{4b^2}(\vec{r}^2 + \vec{r}'^2)}$$

we see that the transition matrix element $|\Sigma \bar{V}_{qq}|^2 \approx (0.02/m_q b^2)^2$ is rather small on the scale of a typical confinement interaction $V_{qq}(r^2=b^2) \sim (1/3m_q b^2)$. Mathematically, we can trace the smallness of the rearrangement potential back to a cancellation between the parts in the transition operator [Eq. (14)] coming from the qq- interaction and the kinetic energy of the six-quarks.

In Fig. 2, we show the real and imaginary part of the local approximation to the separable potential

$$V(\vec{x}) = \int d^3\vec{r} \int d^3\vec{r}' \delta(\vec{x} - \frac{\vec{r} + \vec{r}'}{2}) V(\vec{r}, \vec{r}')$$

as a function of the radial co-ordinate. This function gives an idea of the range and strength of the potential at $E_{\text{rel}} = 0$. Note the local approximation gives a Gaussian shape factor $\exp[-3x^2/2b^2]$ whereas the separable form factors have the form $\exp[-3r^2/4b^2]$ with a larger range.

In order to take into account long range $\bar{N}\bar{N}$ -interactions we have added to this quark rearrangement process a t-channel meson exchange potential, which is deduced by a G-parity transformation from the cut-off Paris potential¹⁵⁾. A cut-off function¹⁶⁾ $f(r) = (r/r_c)^{10}/(1+r/r_c)^{10}$ suppresses the Paris potential for $r \leq r_c = 0.8$ fm, where it is anyhow mostly phenomenological and cannot be G-parity transformed. It should be kept in mind that the t-channel meson exchange process, i.e., processes with two mesons or $q\bar{q}q\bar{q}$ in the s-channel may also have a real short range contribution. In Ref. 17), e.g., a repulsive interaction is found from annihilation into one gluon. To account for these uncertainties, we will present our results for the $\bar{N}\bar{N}$ s-wave cross-sections in three different ways. First we calculate only the effect of the rearrangement potential (R), then we add the known 3S_1 - 3D_1 meson-exchange potential (R+SD) which is mostly due to one pion exchange, and finally we use the full Paris potential (R+SD+C). We always employ the equivalent separable form of the meson exchange potential derived by the Graz group¹⁸⁾. The cut-off potentials for the various (I,S) states are, in general, weakly attractive and go to zero around $r = 0.5$ fm. Inside this range, the form factors oscillate in configuration space due to compensating Yukawa functions. It would be useful to have a less sophisticated

smooth parametrization which is based on a cut-off at larger distances ($r_0 \approx 1$ fm-1.2 fm). In comparison to the dispersive part of the rearrangement potential the meson exchange dominates at distances $r > 0.8$ fm. In Figs. 3a,b, we show the results of a calculation of the inelastic and elastic s-wave cross-sections deduced from $\bar{p}p$ -scattering between 0.18 GeV/c and 0.28 GeV/c. The data are from Ref. 2). The minimal SU4 invariant model of rearrangement with tensor S-D coupling from meson-exchange gives the best description of the data. An additional attraction due to other t-channel meson exchanges increases the elastic cross-section much more than the inelastic cross-section. From the Table, we see, however, that the full Paris potential (R+SD+C) produces a very large charge exchange cross-section which is unrealistic because of missing strength in annihilation. It was already noticed in the framework of empirical optical models¹⁶⁾ that the smallness of the charge-exchange cross-section gives strong constraints on the size and strength of annihilation.

Using the Trueman formula, we can convert the 1S_0 and 3S_1 scattering lengths into energy shifts and widths of the $\bar{p}p$ -atom. We compare these theoretical values with the experimental results $\text{Re}\Delta E_S = (-0.5 \pm 0.3)$ keV and $\Delta E_S = (-0.73 \pm 0.15)$ keV + $i(0.85 \pm 0.39)$ keV from Refs. 19) and 20). In general, the theoretical width of all models is too small, indicating a too weak imaginary strength of the rearrangement potential. The real energy shift has the wrong sign for the otherwise acceptable R+SD model, but adding the full central meson exchange interaction seems to overcompensate the necessary attraction as seen in the Table.

We now summarize. The quark rearrangement model gives a complex optical potential for $N\bar{N}$ -scattering with a form factor of the size of the baryon size and an absolute strength determined by the string tension between quarks, as measured in hadron spectroscopy. This optical potential gives 50% to 90% of the low-energy s-wave inelastic cross-section as recently measured at LEAR. The results depend on how much of the t-channel meson-exchange potential is added. We believe that pion exchange dominates the long range $N\bar{N}$ interaction and, indeed, is necessary to describe the angular distributions. The problem is how to describe unambiguously the matching between the Yukawa picture at long distances and the short range regime presumably dominated by quarks and gluons. Without this problem being solved, we hardly imagine testing further the details of the quark dynamics in $N\bar{N}$ annihilation.

Table

Results of the rearrangement model (R), with tensor meson-exchange potential (R+SD) and with the full separable Graz-form of the Paris potential (R+SD+C) for the s-wave $\bar{p}p$ -interaction ($\sqrt{s}/2 = 50$ MeV). The energy shift ΔE_s is defined by the average of the two spin states. The scattering lengths a are given for $\bar{p}p$ -scattering. (The convention used is $a = -\tan\delta/k$, where k is the relative momentum and δ the phase shift.)

Model	R		R+SD		R+SD+C	
$p_{\text{Lab}} \text{ (MeV/c)}$	150	300	150	300	150	300
$\sigma_{\text{el}} \text{ [mb]}$	9.5	6.6	59.2	42.3	149.6	103.4
$\sigma_{\text{CEX}} \text{ [mb]}$	0.0	0.0	11.7	18.7	45.9	29.1
$\sigma_{\text{ann}} \text{ [mb]}$	80.3	30.9	146.9	39.6	33.7	16.0
$\sigma_{\text{tot}} \text{ [mb]}$	89.8	37.5	217.8	100.6	229.2	148.5
$a_{1S_0} \text{ fm}$	0.06 - i0.31		0.06 - i0.31		1.53 - i0.23	
$a_{3S_1} \text{ fm}$	0.06 - i0.31		-0.70 - i0.26		1.32 - i0.31	
$+\Delta E_s \text{ (keV)}$	-0.05 + i0.27		+0.28 + i0.25		-1.24 + i0.23	

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FIGURE CAPTIONS

Fig. 1: Schematic picture of the quark rearrangement process for $N\bar{N} \rightarrow M_1 M_2 M_3$.

Fig. 2: Local approximation to the rearrangement potential (a) real part, (b) imaginary part. The different curves correspond to various values of Γ .

Fig. 3: Results for the inelastic (a) and elastic (b) s-wave cross-sections of the pure rearrangement model (R), with tensor meson-exchange potential (R+SD) and with the full Paris potential (R+SD+C).

Fig. 1

Fig. 2a

Fig. 2b

Fig. 3