

TEST FOR K^- -HYPERON RELATIVE PARITY

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Numerous experiments which could determine the (K^- , Λ) or (K^- , Σ) relative parities have been suggested in the last several years. In addition to examinations of threshold behaviors, they include experiments with K in flight,¹ reactions on polarized targets,² with polarized beams,^{1,3} and those in which the hyperon emerges in a hyperfragment.⁴ In the latter case, the spin and orbital parity of the fragment present additional unknown quantities.

We wish to discuss here a parity test with stopped K^- utilizing an unpolarized target. To the best of our knowledge, one such test (K^- on d) has been suggested in the literature⁵; however, it requires the assumption that the meson be absorbed predominantly from an s state. A series of papers⁶ has recently appeared giving theoretical discussions in support of this assumption. The arguments leading to the conclusion of pure s absorption involve the strong mixing of the s states into the high- l states through the Stark effect produced by fields of neighboring atoms. It is sufficiently difficult to come to a reliable quantitative conclusion regarding the fraction of p -state capture that we believe that it is useful to present arguments which show that the same kind of experiment suggested by Sirlin and Spitzer,⁵ when applied to He, yields a parity test which is independent of any special assumptions about the K absorption, except of course that it should conserve parity and that the K have zero spin. In this connection, it should be remarked that the dominance of s over p capture is expected to be less in (K , He) than in (K , d).

Since the result has implications for a variety of processes, we prefer to give the discussion in general terms. Consider in the c.m. frame a reaction $K^- + \text{He}^4 \rightarrow$ any number of particles, all with relative momenta which are parallel or antiparallel in a direction which we call \vec{k} , and take parallel to z . Of course, for the K and α one can substitute any two spin-zero particles. The particles in the final state have intrinsic angular momenta which can add up to the total spin values S, S', \dots . We decompose the initial wave function into orbital angular momentum states $|lm\rangle$, but make no assumption here about the relative amplitudes of the latter. The discussion can therefore

be applied to initial plane waves as well as to bound states.

Under a reflection R in the $x-z$ plane (inversion followed by 180° rotation about the y axis) the final state and the initial components transform, respectively, as

$$R |\vec{k} S m_S\rangle = (-1)^{S-m_S} \mathcal{P}_f |\vec{k} S - m_S\rangle,$$

$$R |lm\rangle = (-1)^{l-m} (-1)^l \mathcal{P}_i |l - m\rangle, \quad (1)$$

in which $\mathcal{P}_{f,i}$ represent the products of all intrinsic parities in the final and initial states. Next we make use of the fact that the final-state plane wave contributes no z component of orbital angular momentum; it is at this point in the argument that, in the case of three outgoing particles, we require that the directions of the final momenta coincide. Then angular momentum conservation gives $m_l = m_S = m$. The invariance of the S matrix gives in turn

$$\langle \vec{k} S m | S | lm \rangle = \mathcal{P}_i \mathcal{P}_f (-1)^S \langle \vec{k} S - m | S | l - m \rangle. \quad (2)$$

The result which we shall use follows immediately from Eq. (2):

$$\mathcal{P}_i \mathcal{P}_f (-1)^S = -1 \text{ forbids } m = 0. \quad (3)$$

The experimental implications of this statement must be examined separately in each case. The simplest one, that in which all final particles have zero spin, makes the process completely forbidden if $\mathcal{P}_i \mathcal{P}_f = -1$, in agreement with the result of Bohr.¹ Bohr's symmetry operation can make a stronger statement about this process than does ours, since he can forbid three-particle final states with two arbitrary momentum directions provided that the initial K is in flight in the same plane. His conservation law and relation (3) have a large area of overlap, but they are not identical in content since we have invoked rotational invariance about the z axis in addition to the reflection R .

For the K^- -hyperon parity test, we consider states in which two of the final particles have spin $\frac{1}{2}$ and in which all final-state momenta are

parallel:

$$K^- + \text{He}^4 \rightarrow \left\{ \begin{matrix} \text{H}^3 \\ \text{He}^3 \end{matrix} \right\} + \left\{ \begin{matrix} \Lambda \\ \Sigma \end{matrix} \right\} + n\pi. \quad (4)$$

A specification of the precise nature of the initial state can still be deferred; we denote it by ψ . We now show that a measurement of the sign of the spin correlation averaged over all directions perpendicular to \vec{k} gives a direct determination of $\mathcal{O}_i \mathcal{O}_f$. Indeed, the operator

$$C_{\perp} \equiv \frac{1}{2}(\sigma_x^1 \sigma_x^2 + \sigma_y^1 \sigma_y^2) = \frac{1}{2}(\vec{\sigma}^1 \cdot \vec{\sigma}^2 - \sigma_z^1 \sigma_z^2) \quad (5)$$

is diagonal in S and m . The triplet $m = \pm 1$ states contribute zero to $\langle C_{\perp} \rangle$; the triplet and singlet $m = 0$ states contribute $+1$ and -1 , respectively. Combining this with relation (3), we find the final result

$$\langle C_{\perp} \rangle = -A_{00}(A_{00} + A_{11} + A_{1-1})^{-1}, \quad \mathcal{O}_i \mathcal{O}_f = +1, \quad (6a)$$

$$= +A_{10}(A_{10} + A_{11} + A_{1-1})^{-1}, \quad \mathcal{O}_i \mathcal{O}_f = -1, \quad (6b)$$

in which $A_{Sm} \equiv |\langle \vec{k} S m | s | \psi \rangle|^2$ and can include contributions from all l .

It is clear that $\langle C_{\perp} \rangle$ can be replaced by $\langle C_y \rangle$ if the absorption is from a mesonic atom which has lost all directional memory of its formation process.

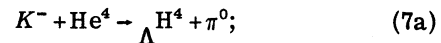
The processes (4) in which no pions emerge have very small branching ratios. When one pion is emitted, it is important to estimate how large an angle θ can be tolerated before condition (6) breaks down, where θ is the c.m. angle between the two momenta which are taken to lie in the x - z plane. To do this, we note that the pion will have momentum $\sim (1 \text{ to } 2)\mu c$. If we assume an interaction radius $\sim \hbar/\mu c$ for (4), pion d waves at the highest will be involved and it is unlikely that $\langle C_{\perp} \rangle$ or $\langle C_y \rangle$ will change sign when $\theta \lesssim \pi/4$. Furthermore, pure s -wave absorption of the K would give $\langle C_y \rangle = \mp 1$ ($\mathcal{O}_i \mathcal{O}_f = \pm 1$) for arbitrary θ . This latter result comes from the Bohr operation, according to which one applies the operator R to initial and final state, but with all spins quantized along y . (It is straightforward to show that $\langle C_{\perp} \rangle$ also carries the unique sign in the s -capture case, but there is no restriction on its absolute value.) Since we can reasonably expect at least a large part of the capture to be from s states, a combination of the above two estimates makes it appear likely that the correlation $\langle C_{\perp} \rangle$ or $\langle C_y \rangle$ resulting from an unpolarized K -mesonic atom will not change sign as θ ranges through all values. A quantitative measurement of $\langle C_y \rangle$ could give a

direct indication of the fraction of the absorption arising from states with $l > 0$,

To know $\langle C \rangle$, one must have simultaneous information on the polarization directions of the nucleus and the hyperon. In special cases knowledge of the individual polarizations can be shown to yield the sign of $\langle C_y \rangle$; the extreme case in which $P_1, P_2 = \pm 1$ is obvious. If the strange particle is a Λ or Σ^+ , its own decay would serve as the polarization signature once the sign of the asymmetry parameter α is settled.⁷ The Σ^0 polarization could also be inferred from that of the daughter Λ using the well-known parity-independent relation⁸ $\langle P_{\Lambda} \rangle = -\frac{1}{3}P_{\Sigma^0}$. The nuclear spin direction affects the right-left asymmetry of the elastic scattering of the recoil nucleus (which has $E_{\text{max}} \sim 50$ Mev, but $\langle E \rangle$ probably much lower) on He, or the asymmetry in a stripping reaction such as $\text{He}^3(\alpha, p)\text{Li}^6$. The dependence of these processes on polarization is apparently not yet known.

We next relate the above discussion to the (K^-, d) atom. It is clear that the preceding arguments rest on the complete correlation between parity and initial angular momentum in Eq. (2); i.e., $l = j$. For the spin-one deuteron this correlation is lost and the general argument no longer holds. If, however, one assumes all processes to come from the s state so that the initial state is pure 1^+ , the correlation is regained. One has again relation (3) and Eqs. (6), all with a change of sign, in agreement with the Sirlin-Spitzer result⁵ when no pions emerge. In the case of pion emission the same argument can be applied if the momenta are parallel, whatever might be the relative orbital angular momentum of the baryons.

We conclude with two examples of other uses of relation (3). The first concerns the reaction



The angular distribution for the products of (7b) when $j = 1$ has been discussed by Dalitz and Downs⁴ using Clebsch-Gordan analysis. For $\mathcal{O}_{\pi} \mathcal{O}_K \mathcal{O}_{\text{Frag}} = +1$, they show that a certain Clebsch-Gordan coefficient vanishes and that therefore the $m = 0$ state is absent, giving rise to a pure $\sin^2 \theta$ distribution. The same result is seen immediately from (3).

Secondly, the relation has a low-energy nuclear physics application as a detector of nuclear alignment. Consider $X(d, \alpha)Y^*$ in which X and Y^* are both 0^+ states. Then, choosing the line of flight

of the deuteron as the z axis, we see that (3) completely forbids the process to proceed from the $m=0$ state for all scattering angles. If this reaction is used to analyze the deuterons coming from some primary reaction, it can serve as a sensitive detector of the angular dependence of $\langle T_{20} \rangle \sim \langle 3S_z^2 - 2 \rangle$.

Although reactions of the above type to low-lying levels of light nuclei are forbidden by isotopic spin conservation, one can break the T selection rule by choosing sufficiently heavy nuclei. An example⁹ is the exothermic reaction $\text{Si}^{28}(d, \alpha)\text{Al}^{26*}$ to the first excited state of Al^{26} .

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$N\bar{N} \rightarrow \pi\pi$ AMPLITUDE AND THE ELECTROMAGNETIC STRUCTURE OF THE NUCLEON*

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Recently, the effect of a pion-pion P -wave resonance on the $N\bar{N}$ to $\pi\pi$ amplitude and the electromagnetic structure of the nucleon has been investigated by Frazer and Fulco.¹ Their calculation suffers from two deficiencies: (1) The "rescattering cut" they employed produces divergent integrals which necessitate a cutoff. (2) Their "one-pole" effective-range formula corresponds to an extremely short-range pion-pion force and is inconsistent with the crossing symmetry of the pion-pion problem.²

The purpose of this Letter is to remedy both of these difficulties. (1) By crossing symmetry, the $N\bar{N}$ to $\pi\pi$ amplitude at zero total energy is identical to the forward pion-nucleon scattering amplitude, which can be calculated directly from experimental information. Having thus determined

the $N\bar{N}$ to $\pi\pi$ amplitude at one point, one can calculate both this amplitude and the nucleon form factors more reliably. In fact, a one-subtraction dispersion relation can be formulated which removes the divergence difficulty. (2) Chew and Mandelstam have observed from the crossing relations for pion-pion scattering that a P -wave resonance produces a long-range repulsive force in addition to the shorter range attraction. We represent such an interaction by a "two-pole" P -wave effective-range formula.

To obtain the normalization of the $N\bar{N} \rightarrow \pi\pi$ amplitude, we start with the pion-nucleon fixed momentum-transfer dispersion relations in the neighborhood of zero momentum transfer. The projection of the $J=1, I=1$ $N\bar{N} \rightarrow \pi\pi$ characteristic amplitudes from the dispersion relations yields

$$f_{+}^{-1}(t) = \frac{1}{p_{-}q_{-}} \left\{ 4f^2 m^3 z_0 Q_1(z_0) + \frac{1}{4\pi^2} \int_{(m+1)^2}^{\infty} ds \left[\frac{p_{-}}{q_{-}} \text{Im}A^{(-)}(s, t) + mz \text{Im}B^{(-)}(s, t) \right] Q_1(z) \right\}, \quad (1)$$

$$f_{-}^{-1}(t) = \frac{\sqrt{2}}{3p_{-}q_{-}} \left\{ 4f^2 m^2 \left[Q_0(z_0) - Q_2(z_0) \right] + \frac{1}{4\pi^2} \int_{(m+1)^2}^{\infty} ds \left[\text{Im}B^{(-)}(s, t) \right] \left[Q_0(z) - Q_2(z) \right] \right\}, \quad (2)$$