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PION PRODUCTION BY TARGET EMISSION

by

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ABSTRACT

The coherent production of pions in a proton nucleus collision is calculated. The energy region considered is below the nucleon-nucleon threshold so that the entire nucleus must take part in the process. The amplitude for this reaction is broken into two parts, the first part being characterized by initial emission of the pion from the projectile and the second part by initial emission from the target nucleus. It is the second part which is emphasised in this talk. The target nuclei ^{12}C and ^{209}Bi are considered. Total production cross sections as a function of energy and some angular distributions are calculated.

INTRODUCTION

Let me begin by trying to put my talk in perspective. To do this I will consider only the one pion rescattering graph. It is known ¹ that a reasonable representation of the $p+p \rightarrow \pi+d$ reaction can be achieved with this type of calculation. This is not to deny the possible importance of ρ -exchange ², interacting Δ 's ³, relativistic effects ⁴ etc. but to try to take the simplest graph for use in the nuclear calculation. Taking a static pion approximation immediately we may write

$$M = -2 \sum_{\vec{q}, i, j} \frac{(\vec{q}' | H_s^i | \vec{q})}{\omega(q)} \frac{(\vec{q} | H_p^j | 0)}{\omega(q)} \quad (1)$$

Here i and j label the nucleons involved in the reaction, $\vec{q}'$ is the final pion momentum and $\vec{q}$ the intermediate pion momentum. H_p^j is the production vertex operator and H_s^i is the corresponding rescattering vertex operator. The nuclear states have been suppressed but are implicit [Eg. $|0\rangle$ means a proton incident on ^{12}C and no pions etc.]. The rescattering operator can be related to the pion nucleon t -matrix and the production vertex is related to γ_5 .

To obtain the two nucleon model ⁵ one calculates first a pion production "potential" by doing the $\vec{q}$ sum (integral):

$$V_{ij} = V(\vec{r}_i - \vec{r}_j) \sim \int d\vec{q} \frac{t(\vec{q}, \vec{q}') \vec{\sigma} \cdot \vec{q}}{\omega^2(q)} e^{i\vec{q} \cdot (\vec{r}_i - \vec{r}_j)} \quad (2)$$

One can then calculate the expectation value in the nucleus of this potential

$$M \sim (\vec{q}' | \sum V_{ij} | 0) \quad (3)$$

This is a perfectly reasonable way to proceed but I shall consider another organization of the work. We may note that in the i, j sum either i or j must be the projectile. If j is the projectile then i just labels the target nucleons. We may recognize

$$\sum_{i=1}^A (\vec{q}' | H_s^i | \vec{q}) \quad (4)$$

as the single scattering approximation to the pion nucleus t -matrix. If we improve the calculation by using the complete multiple scattering t -matrix and add in the plane wave part we recognize this result as

$$M \sim (\phi_\pi(\vec{q}') | H_p^P | 0) \quad (5)$$

which is the usual one nucleon DWIA matrix element.

What about the other part? Let us consider the case where i labels the projectile. In this case the shell structure is more complicated. I will use the ^{12}C target as a specific example. The final state will be some state in ^{13}C and we can write it as an expansion of the final captured nucleon multiplied by core states of ^{12}C

$$|^{13}\text{C}\rangle = \sum C_{\ell j}^{J T E} | J T E \rangle | \ell j \rangle \quad (6)$$

This part of M will now be

$$\sum_{i \in ^{12}\text{C}} \int d\vec{q} \frac{(\vec{q}' | \ell j | H_s^P | \vec{q} \psi_p)}{\omega(q)} (\vec{q} | J T E | H_p^i | 0) C_{\ell j}^{J T E} \quad (7)$$

Since $H_s^i \sim \vec{\sigma} \cdot \vec{q} \tau_i^i$ the core states must be unnatural parity ($0^-, 1^-, 2^-, \dots$) and have isospin one. Note that the matrix element on the left refers only to the projectile and final nucleon and the matrix element on the right contains only target coordinates. In order to make further progress let us expand the ^{12}C core states in particle-hole promotions from the ^{12}C ground state i.e.

$$| J T E \rangle = \sum B_{\ell_h j_h \ell_p j_p}^{J T E} \left\{ | \ell_h j_h \rangle | \ell_p j_p \rangle \right\}^{J T} \quad (8)$$

Now the production matrix element can be written as

$$(\bar{q} J T E | H_p^j | 0) \sim \sum_{\ell_h j_h \ell_p j_p} B_{\ell_h j_h \ell_p j_p}^{J T E} \frac{1}{\sqrt{\omega(q)}} Y_J(\hat{q}) Q_{ph}^{JL}(q) a_{JL} \quad (9)$$

$$Q_{ph}^{JL}(q) = \alpha f(q^2) \left(\frac{q}{m_\pi}\right) \hat{c}_{ph}^{JL} \int_0^\infty dr r^2 j_L(qr) R_p(r) R_h(r) \quad (10)$$

where $\alpha=2$ for nucleons, $\alpha = 8/3$ for Δ -isobars. This last factor of α takes account of the fact that we could also excite a Δ -hole component. The reason for writing the function in this way is to make connection with the work of Toki and Weise, ⁵ and the Lyon group ⁷ as well. Toki and Weise show how the Q function may be corrected by means of RPA to include the effects of multiple particle-hole and delta-hole states. If we represent the simplest form of Q as defined above as

then, taking into account the possibility of exciting virtual particle-hole states by either pion exchange or heavy meson exchange, the Q gets replaced by $\tilde{Q}$ where $\tilde{Q}$ can be represented by a sum of diagrams.

Note that, depending on one's point of view, this correction can be regarded as a correction to the basic particle-hole wave function or a modification to the pion propagator. The main effect in any case is to give more support to the high momentum components. The rest of the nucleons in the nucleus "back up" the simple particle-hole wave function so that the pion-nucleus vertex is given enhancement from the coherent many-body effects. This actually provides a more-than-two nucleon theory of pion production.

Leaving off the spin and isospin indicies one can write a typical element in the sum for M as the deceptively simple expression

$$\int d\vec{r} \phi_{\pi}^{(-)*}(\vec{r}) \psi_f(\vec{r}) f(\vec{q}, \vec{q}', \vec{p}, \vec{p}') A(\vec{r}) \psi_i(\vec{r}) \quad (11)$$

where

$\psi_i(\vec{r})$ is the initial proton wave function,

$A(\vec{r})$ is the function $Q_{ph}^{JL}(q)$ in r-space,

$f(\vec{q}, \vec{q}', \vec{p}, \vec{p}')$ is a function derived from the pion-nucleon t-matrix as an operator in all four momenta,

$\psi_f(\vec{r})$ is the final bound nucleon wave function,

$\phi_{\pi}^{(-)}(\vec{r})$ is the final pion wave function.

This has the form of a DWIA matrix element for knock out of a pion followed by a capture of the incident nucleon. The reason that the computation of this integral is not simple is the presence of the operators $\vec{q}, \vec{q}', \vec{p}, \vec{p}'$ in f . Here $\vec{q}$ is the gradient operator on $A(\vec{r})$, $\vec{q}'$ is the gradient operator on $\phi_{\pi}^{(-)}(\vec{r})$, $\vec{p}$ is the gradient operator on $\psi_i(\vec{r})$ and $\vec{p}'$ is the gradient operator on $\psi_f(\vec{r})$. It is only the assumption of a fairly simply form for f that makes computation possible. The expression used for f is

$$f(\vec{t}, \vec{t}') = (\lambda_0 + \lambda_1 \vec{t} \cdot \vec{t}' + \tilde{\lambda}_1 \vec{\sigma} \cdot \vec{t} \times \vec{t}') v(t)v(t') \quad (12)$$

where

$$\vec{t} \equiv \vec{q} - \frac{\mu}{m} \vec{p} \quad \text{and} \quad \vec{t}' \equiv \vec{q}' - \frac{\mu}{m} \vec{p}'$$

A schematic representation of this integral is given in fig. 1. The initial proton interaction is a standard (non-relativistic) proton nucleus optical potential.⁸ The pion "wave function" was calculated⁹ with the methods of Toki and Weise. The needed shell model coefficients were taken from the Los Alamos version of Glasgow code¹⁰ as implemented by Haxton and Dubach.¹¹ The final nucleon bound state wave functions were solutions in a Saxon-Woods potential. The shell model coefficients were obtained as above. The final pion wave function was taken from the coordinate space non-local optical model code developed in T-5 at Los Alamos.¹²

Figure 1. A schematic representation of expression 11.

I am going to consider only the target emission in most of what follows. The reasons for this are several. The first motivation is to separate the problem into manageable pieces and to study them independently. Since π^- production may well be dominated by target emission there is a chance to look at it (almost) alone. Another reason is the hope that the sensitivity to the final optical potential will be less than in the projectile emission since more vertices are present in the target emission. Another sensitivity in the projectile emission is due to the treatment of true pion absorption with an incident off-shell pion. Presumably the possibility of reabsorption is much greater since the momentum of the pion is much nearer the pole for true absorption in the optical potential.¹³ This would affect the projectile emission more than the target emission because in the latter the pion is more likely to be on-shell when it leaves the π -nucleon t-matrix.

II. ENERGY DEPENDENCE FOR $^{12}\text{C}(p, \pi^+)^{13}\text{C}$

There are many possible calculations that could be done with this technique. I will try to focus primarily on a single simple point, namely the energy dependence of the total pion production cross section. This data is now available on a large number of nuclei but for the first comparison with experiment I shall restrict myself to several states of ^{13}C . It was noted by the Indiana group¹⁴ that there seemed to be two types of energy dependence of the total cross section; one in which there was an initial rapid rise with a "bend" to give a much more slowly rising function and a second type which simply increased from zero with no abrupt change in slope. The first type of behavior seemed to be associated with single particle states and the second type with 2p1h states. An exception to this rule is the $9/2^+$ state at 9.5 MeV.

Since simple projectile emission populates single particle states only, while target emission populates single particle and 2plh states we might conjecture that the separation of reaction mechanisms might be made in this way.

The first step in this investigation was the calculation of a total cross section with projectile emission, figure 2 shows the results of a calculation using the Keating-Wills-Tsangarede, 15 code.

Fig. 2. One-nucleon-model calculations showing the dependence on the non-locality of the optical potential. Note that this non-locality is a total effective value including e.g. the motion of the Δ .

The curve labeled "Lap" was calculated using the damped Laplacian form. The other curves were calculated using the T-5 non-local code with the off-shell range denoted as α . One might expect a similar variation with other optical model parameters, such as p^2 terms.

We may see that (within limits) the very low energy part is little affected by the optical potential and the variation is largely controlled by phase space and Coulomb penetrability. We note that a change in slope can be obtained at about the right energy. Thus the hoped-for separation hypotheses is still alive although it is somewhat disturbing that the calculated cross section showing the desired shape is so far below the data. The uncertainties in the optical potential may allow us to hope for the correction of this feature.

The next step is to see if the target emission can explain pion production to the 2plh states. Let us first look at some characteristics of the target production mechanism. The effect of each of the terms in Eq. 12 is shown individually in Fig. 3.

Fig. 3. Characteristic shapes of a target emission calculation showing the dependence on the various parts of the π -nucleon t -matrix.

The curve labeled "1" contains only the S-wave pion-nucleon scattering and thus shows the form factor and distortion effects only. The other two curves show the $\cos^2\theta$ and $\sin^2\theta$ dependences as well. These features are very clear in the case presented, but not quite so clear in others. The spin dependent term is too large for the ^{13}C states. It is not known whether this is due to the fact that the pion-nucleon t-matrix should be modified in the nuclear medium to reduce the spin dependent piece, whether some feature of the calculation should be modified to reduce this term or if we are simply barking up the wrong graph. In any case, for the remainder of the comparisons with ^{13}C I shall set this term to zero.

The calculations are classified to some extent by the core excited state. For p-wave capture in the final state the positive parity states will be described by 0^+ or 2^+ (or higher in some cases) core states.

Fig. 4. The hole particle components entering into the structure of the 2^- core state.

The hole-particle pairs allowed for 0^- are $1s_{\frac{1}{2}}^{-1} 1p_{\frac{1}{2}}, 1p_{\frac{1}{2}}^{-1} 2s_{\frac{1}{2}}$ and $1p_{\frac{3}{2}}^{-1} 1d_{\frac{3}{2}}$. For 2^- the available states

are $1s_{\frac{1}{2}}^{-1} 1p_{\frac{3}{2}}, 1p_{\frac{3}{2}}^{-1} s_{\frac{1}{2}}, 1p_{\frac{1}{2}}^{-1} 1d_{\frac{3}{2}}, 1p_{\frac{1}{2}}^{-1} 1d_{\frac{5}{2}}, 1p_{\frac{3}{2}}^{-1} 1d_{\frac{3}{2}}, 1p_{\frac{3}{2}}^{-1} 1d_{\frac{5}{2}}$

The hole-particle A functions for the 2^- case are shown in figure 4.

The actual $A(r)$ used in Eq. 11 is a linear combination of these functions with coefficients given by the shell model. It is common for one or two of these basis functions to dominate. The shape of $A(r)$ has a large effect on modification of the angular distribution.

Fig. 5. The effect of variation of some of the input parameters on the angular distribution. The solid ("std") curve is calculated for c (the W-S radius for ^{12}C in the optical potential) $\approx 2.41\text{fm}$, $b_0 = (2.7, .9)$, $b_1 = (7.8, .5)$, r_b (radius of the W-S well for computing the final bound state) $= 1.31\text{ fm}$, and no enhancement from the RPA corrections.

The first state considered is the $6.86 \frac{5+}{2}$ state in ^{13}C . This state is dominated by components

$$\left\{ \left[\begin{matrix} 1p^{-1} & 2s \\ \frac{3}{2} & \frac{1}{2} \end{matrix} \right]_2 \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}} \quad \left\{ \left[\begin{matrix} 1p^{-1} & 1d \\ \frac{3}{2} & \frac{5}{2} \end{matrix} \right]_2 \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}}$$

and

$$\left\{ \left[\begin{matrix} 1p^{-1} & 1d \\ \frac{1}{2} & \frac{5}{2} \end{matrix} \right]_2 \times 1p_{\frac{3}{2}} \right\}_{\frac{5}{2}}.$$

The $1p1/2$ capture dominates and I show only these results in Fig. 5, i.e. all six components are included in the 2^- core but only the $p1/2$ coupling is calculated. Figure 5 shows some angular distributions compared with data at 185 MeV. As can be seen the qualitative agreement is quite good with only moderate variations in shape corresponding to rather substantial modifications in the model.

Figure 6 shows the angular distribution as a function of energy. Detailed agreement leaves a great deal to be desired but the dependence of magnitude on energy is good.

Fig. 6. Energy dependence of the angular distributions for the 6.86 MeV state corresponding to the "std" curve in Figure 5. The dash-dot curve is the same as in Figure 5.

Figure 7. The energy dependence of the integrated cross section for the 6.86 MeV state.

This is reflected in Fig. 7 where the integrated cross section is plotted. Thus we see that the gradual energy dependence is correctly given by this calculation.

The third step in the study of the separation hypothesis requires that the target emission calculation fail to give the energy dependence of a single particle state. For this case I have chosen the 3.09 $1/2^+$ final state in ^{13}C . This state is well represented by $^{12}\text{C} \times 2s_{1/2}$. The target emission must destroy the core, however, so to restore the state the captured particle must fill the hole. Thus since $1/2^+$ can be obtained by $0^- \times 1p_{1/2}$ or $2^- \times 1p_{3/2}$ the dominant components are

$$\left\{ \left[\begin{array}{cc} 1p_{1/2}^{-1} & 2s_{1/2} \\ 2 & 2 \end{array} \right]_0 \times 1p_{1/2} \right\}_{1/2}$$

and

$$\left\{ \left[\begin{matrix} 1p_3^{-1} & 2s_1 \\ \frac{1}{2} & \frac{1}{2} \end{matrix} \right]_2 \times 1p_3 \right\}_{\frac{1}{2}}$$

Figure 8 shows the result of the calculation from s- and d-core states (all hole-particle states included).

Figure 8. Separate angular distributions for the 3.09 MeV state.

Note that at high energy where the experimental minimum is deep the minima of the two pieces nearly coincide. At low energy where the experimental angular distribution becomes much flatter then the two theoretical minima do not agree. It looks almost accidental that the d-wave comes up in the back angles but recall that the relative strengths will be given by angular momentum coefficients obtained by the decomposition of the single particle state.

In figure 9 the cross sections from the coherent sum of the s- and d- state core are shown. The general agreement with the angular distributions, in particular the energy dependence of the forward hemisphere cross sections, is largely satisfactory. In fact the agreement seems rather better for this single particle state than for the 2plh state.

Figure 9. Combined angular distributions for the 3.09 MeV state.

The corresponding integrated cross sections are shown in figure 10.

Figure 10. Separate and combined integrated cross sections for the 3.09 MeV state.

We note that the s- and d- state energy dependences are different with the s- state showing a dependence like the single particle shape. The s- and d- states tend to match better at low energy and the coherence holds up the cross section there.

Thus we see that a separation of diagrams by the energy dependence is not possible. This does not mean that a separation by structure cannot be made, just that the association of a dominance of target emission with 2plh states is not true.

A fairly typical asymmetry is shown in fig. 11. This is in reasonable qualitative agreement with recent data.¹⁶

Figure 11. An example of asymmetry for the 3.09 MeV state. Only a 0^+ core is included.

I shall briefly consider one other state, the 3.85 state in ^{13}C . Since this is also a $\frac{5}{2}^+$ state like the 6.86 one might expect the pion production to be similar. However, the particle-hole composition is rather different. For this state the dominant (predicted) contributors are

$$\left\{ \left[\begin{matrix} 1p_{3/2}^{-1} & 1d_{5/2} \\ 2 & 2 \end{matrix} \right]_2 \times 1p_{3/2} \right\} \frac{5}{2}$$

and the $p_{\frac{1}{2}}$ couplings

$$\left\{ \left[\begin{matrix} 1p_{\frac{1}{2}}^{-1} & 1d_{\frac{5}{2}} \\ 2 & 2 \end{matrix} \right] \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}} \quad \left\{ \left[\begin{matrix} 1p_{\frac{3}{2}}^{-1} & 1d_{\frac{5}{2}} \\ 2 & 2 \end{matrix} \right] \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}}$$

$$\left\{ \left[\begin{matrix} 1p_{\frac{3}{2}}^{-1} & 2s_{\frac{1}{2}} \\ 2 & 2 \end{matrix} \right] \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}} \quad \left\{ \left[\begin{matrix} 1p_{\frac{3}{2}}^{-1} & 1d_{\frac{3}{2}} \\ 2 & 2 \end{matrix} \right] \times 1p_{\frac{1}{2}} \right\}_{\frac{5}{2}}$$

The characteristics of the measured angular distributions are rather similar to the other $\frac{5}{2}^+$ state at 6.86. According to the shell model the underlying structure is different. This difference is reflected in the angular distributions of figure 12. I only give the separate pieces but it seems clear that, while the general magnitude may not be too bad there is a great deal lacking in further agreement. This leads us to think that a fair sensitivity to details of nuclear structure exists.

Figure. 12. Separate angular distributions for the 3.85 MeV state.

III. ENERGY DEPENDENCE IN $^{209}\text{Bi}(p,\pi^-)^{210}\text{At}$

Let me now consider a second study of the energy dependence of the target emission. A large part of negative pion production may be reasonably thought of as coming from target emission. The projectile emission may produce π^- by π^0 production followed by charge exchange and more exotic processes (such as Δ^+ mechanisms) can exist as well. I shall neglect these possibilities and assume that the negative pion is produced from one of the neutrons in the target so that target emission is the total mechanism for negative pion production, admittedly an extreme point of view.

The reaction considered is $^{209}\text{Bi}(p,\pi^-)^{210}\text{At}$ followed by the evaporation of neutrons. From the residual isotopic distribution, information about the original energy spectrum can be inferred. To make this connection some additional assumptions will be needed. First we shall assume that the cross section to a state at a given energy can be represented (on average) by the cross section to a "typical" state.

For simplicity this state is taken to be a pure configuration.

$$\left\{ \begin{bmatrix} 1h_9^{-1} & 2g_9 \\ \frac{2}{2} & \frac{2}{2} \end{bmatrix} \times \frac{2g_9}{2} \right\} \frac{9}{2}$$

The cross sections calculated under this assumption are shown in figure 13. There are several interesting features in these curves.

Figure 13. Calculated energy dependence to the "typical" state in the $^{209}\text{Bi}(p,\pi^-)^{210}\text{At}$ reaction.

A rough idea of the isotopic abundance of ^{210}At to be left after this process can be estimated by integrating the first 10 MeV of this function (it should be multiplied by a level density, and will be shortly). We note that this integral rises from threshold and then drops quickly, peaking at the surprisingly low energy of 10-20 MeV above threshold. This feature has already been noted both experimentally and theoretically.¹⁷ Another feature of interest is the importance of the energy dependence of the pion-nucleus optical potential. The dotted curves show the energy variation of these basic cross sections fixing the optical parameters at their value at 50 MeV. The solid curves were calculated with all parameters kept fixed except for the imaginary part of b_1 which was taken as a linear combination of a constant plus the "free" value such that the elastic scattering is well represented at 50 MeV.

To calculate the spectral strength we must know the density of our "typical" states. In a simple model of equal spacing of single particle states, for high energies where Pauli exclusion effects should be unimportant, the level density should vary as the energy raised to the number of "particles" minus one. Thus, for these 2p1b states, the density should vary at E^2 . To make the density a little more realistic at low energies the form

$$\rho(E) = .004 (E+15)^2$$

(E in MeV) was used. The spectrum calculated using this function is shown in figure 14.

Figure 14. Calculated spectrum for negative pion production from protons on ^{209}Bi . The curves labeled E^1 and E^3 are shown to estimate sensitivity to the level density assumption.

Also shown is the shape of the spectrum for two neighboring powers (normalized to the same peak value). Since the normalization of the level density and the spectroscopic factors are very crude the results can be considered as normalized to the total integrated data to all observed isotopes.

Following an evaporation of successive neutrons, the original ^{210}At spectrum can be converted to a cross section for production of various isotopes. The comparison with data ¹⁸ is shown in figure 15.

Figure 15. Comparison of the calculated relative isotope production cross section (open circles) with the data of Ref. 18.

The agreement is very good. The production of the various isotopes is extremely sensitive to the bombarding energy, especially near threshold. These cross sections are shown in figure 16.

Figure 16. The dependence on proton bombarding energy of the cross section for production of At isotopes from a Bi target.

Clearly this technique provides an alternative to the use of Fermi motion for inclusive sub-threshold production of mesons. More detailed comparison with data is needed to select between these two models.

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