Phenomenological model for inelastic scattering of 800 MeV/c pions from nuclei

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We present the results of our calculations for the inelastic scattering of 800 MeV/c charged pions from the ^{12}C nucleus within a strong absorption framework in the adiabatic approximation. For this purpose we extend and generalize our previous model which has been successfully applied to the description of pion-nucleus elastic scattering at 800 MeV/c. This has been achieved by introducing a phenomenological approach involving complex angular momenta in the spirit of Regge. The results of the present calculation are found to be in very good agreement with the experimental data. We also discuss the validity of this approximation within the context of high energy pion-nucleus inelastic scattering and the implications of the results obtained regarding the structure of the 2^+ (4.44 MeV) and 3^- (9.64 MeV) states of ^{12}C.

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I. INTRODUCTION

There has been much work done relating to pion-nucleus scattering in the energy region around the (3,3) resonance and at lower energies [1]. At higher energies, where pion laboratory kinetic energies T_π exceed about 400 MeV, there have been comparatively few studies. One major study in this high energy regime involves very accurate measurements made at Brookhaven National Laboratory (BNL) [2] of the differential cross sections of charged pions at $T_\pi \approx 673$ MeV elastically scattered from ^{12}C and ^{40}Ca and inelastically scattered from ^{12}C. Most of the analyses of this work have been confined to the elastic scattering results [2–5]. With regard to the inelastic scattering data, analyses have been performed (1) within the framework of the distorted wave Born approximation [2] and the distorted wave impulse approximation in conjunction with an eikonal approximation [6] and (2) within the framework of Glauber theory [7]. In most of these studies the results show adequate overall agreement with experiment.

In the present work we take an approach to the high energy inelastic scattering based on a strong absorption model. The success of our previous study [8] within this framework in describing the main features of the pion-nucleus elastic scattering results at 673 MeV has motivated our further development of the phenomenological model to describe the inelastic scattering results. In addition, the model to be described in the present work represents a further extension of an earlier approach which was formulated within the framework of partial wave expansions in the adiabatic approximation in the spirit of Chou and Yang [9]. The model has given very successful descriptions [10,11] of (1) the inelastic scattering results obtained at Saclay of 1.37 GeV α particles involving the low lying excited states of several isotopes of Ca and (2) the inelastic scattering results obtained at CERN of 180 MeV antiprotons involving the low lying excited states of ^{12}C and ^{18}O. In further developing the model in the present study we have specifically taken into account the enhanced transparency of the nucleus to the 673 MeV kinetic energy pions as we have done in our earlier investigation of the elastic scattering differential cross sections.

In order to describe the high energy inelastic scattering of pions from nuclei we exploit the short wavelength of the incident pions, as we have done for the elastic case, to extract further information on the geometry of the target nucleus, that is, an effective nuclear radius R_0 and effective surface thickness a. These properties in turn allow us to infer the distribution of the nuclear matter density. In addition we account for the enhanced transparency of nuclear matter to the incoming pions by introducing an imaginary component to the S matrix. We accomplish this within the framework of complex angular momenta which is very much in the spirit of the approach to scattering theory first introduced by Regge [12].

The main features of the phenomenological model are presented in Sec. II. The results of our calculations and their comparison to experimental data are given in Sec. III. Finally, we summarize our results and give conclusions in Sec. IV.

II. ASSUMPTIONS AND FORMALISM OF THE MODEL

Since the assumptions and details of the model can be found in Refs. [8,10,11] we concentrate only on those aspects that are new in the present formulation (we follow the notations and definitions of these references). The starting point is the general expression for the scattering amplitude for a spinless meson incident on a spinless nucleus given in terms of its usual partial wave expansion

$$f(\theta) = \frac{i}{2k}\sum_{l=0}^{\infty}(2l+1)[1-S(l)]P_l(\cos\theta), \tag{1}$$

where k is the pion wave number in the center of mass system with the S matrix $S(l) = \exp(2i\delta_l)$ given in terms of the nuclear phase shifts δ_l. In the previous studies [10,11] of scattering of antiprotons and α particles from nuclei the S matrix was taken as purely real and of a functional form which depends only on the impact parameter b in combination with the nuclear radius R and effective surface thickness a. Due to the presence of strong interaction and high matter density in the interior of the nucleus, for very small values of the impact parameter the nucleus is assumed to be approximately black to the high energy strongly interacting particles. Consequently, contributions to the elastic scattering and inelastic scattering with the excitation of the low lying collective states may be assumed to be mostly due to the interaction taking place in a finite region of the nuclear surface. Therefore the specific functional form for the S matrix is taken to be

$$S(b) = \left[1 + \exp\left(\frac{R-b}{a}\right)\right]^{-1}. \tag{2}$$

In order to introduce an imaginary component to the S matrix it is convenient to express it in angular momentum representation as

$$S(l) = \left[1 + \exp\left(\frac{L_R - l}{\Delta}\right)\right]^{-1}, \tag{3}$$

where $\Delta = ak$, $b = (l + 1/2)/k$, and $R = (L_R + 1/2)/k$. To account for the partial transparency of the nucleus to the incoming pions we introduce an imaginary component to the S matrix. An extension of the form given by Eq. (3) to include an imaginary component is most naturally made within the framework of complex angular momenta, z, following Regge [12]. In the present case we exploit the fact that $\{1 + \exp[(L_R - z)/\Delta]\}^{-1}$ is an analytic function of the complex variable z and take S at the outset as being defined for complex angular momenta. We therefore take the following form for the S matrix:

$$S(l + i\mu) = \left[1 + \exp\left(\frac{L_R - l - i\mu}{\Delta}\right)\right]^{-1}, \tag{4}$$

where μ is a real number, the value of which controls the contributions to the scattering amplitude from the real part of the nuclear phase shifts. Since $S(z)$ is an analytic function we can expand Eq. (4) about the point l on the real axis in a convergent series in the quantity μ/Δ. At high energies we expect this quantity to be small (indeed the unitarity of the S matrix requires that $|\mu/\Delta| \leq \pi/2$). If we take only the first order term in μ we get the form of the S matrix used in our study of pion-nucleus elastic scattering [8] at 800 MeV/c where typical values of μ/Δ were in the range of ≈ 0.17 to ≈ 0.38.

The above formulation of the problem allows us to follow a consistent approximation procedure and obtain analytic closed form expressions for the scattering amplitude. We substitute the S matrix from Eq. (4) into the general expression Eq. (1) and use the relation

$$\sum_{l=0}^{\infty}(2l+1)P_l(\cos\theta)\exp(-\alpha|l+1/2|) = \frac{2\sinh\alpha}{(2\cosh\alpha - 2\cos\theta)^{3/2}} \text{ for } \alpha > 0 \tag{5}$$

to write the result in the form

$$f(\theta) = \lim_{\alpha\to 0}\frac{1}{ik}\sum_{l=0}^{\infty}(l+1/2)S(l+1/2+i\mu)P_l(\cos\theta)\exp(-\alpha|l+1/2|), \tag{6}$$

where we will be writing the argument of S in the form $(T - t - i\mu)/\Delta$ with $t = l + 1/2$, $T = L_R + 1/2$. We Taylor expand S about the point t and make use of the relation

$$\frac{d^n S(t)}{dt^n} = (-1)^n \frac{d^n S(t)}{dT^n} \tag{7}$$

which leads to the expression

$$f(\theta) = \sum_{n=0}^{\infty}\frac{(-i\mu)^n}{n!}\frac{d^n}{dT^n}\left[\lim_{\alpha\to 0}\frac{1}{ik}\sum_{l=0}^{\infty}(l+1/2)S(t)P_l(\cos\theta)\exp(-\alpha|l+1/2|)\right]. \tag{8}$$

The quantity in brackets has been evaluated in Ref. [10] and is given by $ig(\theta,\Delta)TJ_1(T\theta)/(\theta k)$ up to the term of order Δ/T which is assumed to be a small quantity. The form factor $g(\theta,\Delta)$ is given by

$$g(\theta,\Delta) = \left(\frac{\theta}{\sin\theta}\right)^{1/2}\frac{\pi\Delta\theta}{\sinh(\pi\Delta\theta)}. \tag{9}$$

Upon substitution into Eq. (8) we are led to the final result for the scattering amplitude

$$f(\theta) = ig(\theta,\Delta)\frac{(kR - i\mu)}{k\theta}J_1[(kR - i\mu)\theta], \tag{10}$$

where we have replaced T by kR.

In order to fully characterize both the elastic and the inelastic scattering we need to introduce the appropriate dynamical model of the target nucleus and evaluate the general expression Eq. (10) for the scattering amplitude. In the present case we accomplish this within the framework of the collective model [13,14] by describing the nuclear surface in terms of spherical harmonics as follows:

$$R(\Theta,\Phi)=R_0\left[1+\sum_{\lambda,\nu}\alpha_{\lambda\nu}(t)Y^*_{\lambda\nu}(\Theta,\Phi)\right], \tag{11}$$

where R_0 is the effective equilibrium radius of the nucleus and the $\alpha_{\lambda\nu}(t)$ are the time dependent collective coordinates of the nuclear surface motion. We substitute Eq. (11) into the general expression for the scattering amplitude, Eq. (10), in order to obtain an expression in terms of the collective coordinates. To accomplish this we first write the Bessel function in Eq. (10) in terms of its integral representation

$$zJ_1(z)=\frac{1}{2\pi}\int_0^z dz'z'\int_0^{2\pi}d\Phi\ \exp(-iz'\cos\Phi), \tag{12}$$

where $z=(kR-i\mu)\theta$. We write Eq. (12) as the sum of two integrals, one from 0 to $z_0=(kR_0-i\mu)\theta$ and the other from z_0 to z. The first integral is easily evaluated and by change of variable we express the second integral as a two dimensional integral in the plane normal to the direction of the incident beam. The scattering amplitude is then written as

$$f(\theta,\Delta,R_0,\mu;\alpha)=ig(\theta,\Delta)\left[\frac{(kR_0-i\mu)}{k\theta}J_1[(kR_0-i\mu)\theta]+I(\theta,R_0,\mu;\alpha)\right], \tag{13}$$

where

$$I(\theta,R_0,\mu;\alpha)=\frac{1}{2\pi\theta^2k}\int_{x_0}^{x}dx'(x'-i\mu\theta)\int_0^{2\pi}d\Phi\ \exp[-i(x'-i\mu\theta)\cos\Phi] \tag{14}$$

and $x=kR\theta$, $x_0=kR_0\theta$. We now make use of the fact that the diffraction of incoming waves by a strongly absorptive, nearly spherical object, is equivalent, with respect to the shadow forming wave, to the diffraction by the object's projection onto the area of integration. In the present case it is then correct to first order in the collective coordinates to evaluate the surface deformations of the target nucleus in the plane normal to the direction of the incoming beam. We therefore take the value of the spherical harmonics at polar angle $\pi/2$ in Eq. (11) for the nuclear surface coordinate. We calculate the x' integral in Eq. (14) first and then substitute Eq. (11) into the result retaining only terms which are linear in the collective coordinates. In this way we then obtain

$$I(\theta,R_0,\mu;\alpha)=\frac{(kR_0-i\mu)}{2\pi}R_0\sum_{\lambda+\nu=\text{even}}\alpha_{\lambda\nu}(t)\beta(\lambda,\nu)(-1)^{(\lambda+\nu)/2}\int_0^{2\pi}d\Phi\ \exp(-iz_0\cos\Phi),\exp(-i\nu\Phi)\,, \tag{15}$$

where

$$\beta(\lambda,\nu)=\left(\frac{2\lambda+1}{4\pi}\right)^{1/2}\frac{[(\lambda-\nu)!(\lambda+\nu)!]^{1/2}}{[(\lambda-\nu)!!(\lambda+\nu)!!]}. \tag{16}$$

To complete the calculation we note that only the real part of $\exp(-i\nu\Phi)$ will contribute to the integral over Φ in Eq. (15) which is then simply a defining relation for Bessel functions

$$2\pi(-i)^{|\nu|}J_{|\nu|}(\rho)=\int_0^{2\pi}d\Phi\ \exp(-i\rho\cos\Phi)\cos(\nu\Phi) \tag{17}$$

for ρ complex. Substitution of Eq. (15) into Eq. (13) by virtue of Eq. (17) yields the following result for the scattering amplitude:

$$f(\theta,\Delta,R_0,\mu;\alpha)=ig(\theta,\Delta)(kR_0-i\mu)\left\{\frac{J_1[(kR_0-i\mu)\theta]}{k\theta}\right.$$
$$\left.+R_0\sum_{\lambda+\nu=\text{even}}\alpha_{\lambda\nu}(t)\beta(\lambda,\nu)(-1)^{(\lambda+\nu)/2}(-i)^{|\nu|}J_{|\nu|}[(kR_0-i\mu)\theta]\right\}. \tag{18}$$

For the case of a purely real S matrix corresponding to $\mu=0$ the general result given by Eq. (18) for the scattering amplitude reduces to the result previously obtained for the case of scattering from nuclei of antiprotons and α particles [e.g., see Eq. (25) of Ref. [10]].

In the present investigation of pion-nucleus inelastic scattering we take the $\alpha_{\lambda\nu}(t)$ as the collective coordinates of nuclear surface oscillations [10]. Upon second quantization, the $\alpha_{\lambda\nu}(t)$ are then interpreted as operators expressed in terms of boson creation and annihilation operators $b^\dagger_{\lambda\nu}$ and $b_{\lambda\nu}$, respectively, for phonons with angular momentum quantum numbers λ and ν. Therefore we have

$$\alpha_{\lambda\nu}(t)=\left(\frac{\hbar\omega_\lambda}{2C_\lambda}\right)^{1/2}[b_{\lambda\nu}+(-1)^\nu b^\dagger_{\lambda(-\nu)}], \tag{19}$$

where the ω_λ are the oscillator frequencies and C_λ the oscillator restoring force parameters. The scattering amplitude Eq. (18), being linear in the collective coordinates, is then seen to involve single phonon excitations.

We now calculate the angular distributions of the elastic and inelastic scattering cross sections by forming the matrix elements for the operators $\alpha_{\lambda\nu}(t)$ in Eq. (18) between the ground state $|0\rangle$ and the single phonon excited states $|1;LM\rangle$ where L is the angular momentum of the phonon and M is its projection along the z axis. The elastic scattering amplitude is given by

$$f_{\rm el}(\theta) = \langle 0|f(\theta,\Delta,R_0,\mu;\alpha)|0\rangle = ig(\theta,\Delta)\frac{(kR_0-i\mu)}{k\theta}J_1[(kR_0-i\mu)\theta]. \quad (20)$$

If we expand our general expression Eq. (20) for the elastic scattering amplitude as a power series in μ and retain only terms up to first order we obtain the form of the amplitude utilized in our previous investigation of pion-nucleus elastic scattering [8]. We have in fact determined by calculation that owing to the smallness of the parameters μ obtained in that study, the contribution of the higher order terms to the elastic scattering amplitude is negligible for the cases considered there.

The matrix elements for the operators $\alpha_{\lambda\nu}(t)$ between the ground state and the single phonon state $|1;LM\rangle$ are easily evaluated from which the inelastic scattering amplitude immediately follows:

$$f_{\rm vib}(\theta;0^+\to 1,LM) = ig(\theta,\Delta)R_0(kR_0-i\mu)\left(\frac{\hbar\omega_L}{2C_L}\right)^{1/2}\beta(L,M)(-1)^{(L+M)/2}(-i)^{|M|}J_{|M|}[(kR_0-i\mu)\theta], \quad \text{for } L+M=\text{even}\,.$$
$$= 0, \quad \text{for } L+M=\text{odd}. \quad (21)$$

The differential cross sections for the inelastic scattering are then obtained by taking the complex square of Eq. (21) and summing over the magnetic quantum number M:

$$\frac{d\sigma}{d\Omega}(0^+\to L) = [g(\theta,\Delta)]^2R_0^2[(kR_0)^2+\mu^2]\left(\frac{\hbar\omega_L}{2C_L}\right)\left(\frac{2L+1}{4\pi}\right) \times \sum_{\substack{M=-L\\ L+M=\text{even}}}^{L}\frac{[(L-M)!(L+M)!]}{[(L-M)!!(L+M)!!]^2}|J_{|M|}[(kR_0-i\mu)\theta]|^2. \quad (22)$$

We readily obtain the differential cross sections of inelastic scattering with excitation of one phonon states of the two multipolarities that are relevant to the present study:

$$\frac{d\sigma}{d\Omega}(0^+\to 2^+) = [g(\theta,\Delta)]^2R_0^2[(kR_0)^2+\mu^2]\left(\frac{5}{4\pi}\right)\left(\frac{\hbar\omega_2}{2C_2}\right) \times\left[\frac{1}{4}|J_0[(kR_0-i\mu)\theta]|^2+\frac{3}{4}|J_2[(kR_0-i\mu)\theta]|^2\right], \quad (23)$$

$$\frac{d\sigma}{d\Omega}(0^+\to 3^-) = [g(\theta,\Delta)]^2R_0^2[(kR_0)^2+\mu^2]\left(\frac{7}{4\pi}\right)\left(\frac{\hbar\omega_3}{2C_3}\right) \times\left[\frac{3}{8}|J_1[(kR_0-i\mu)\theta]|^2+\frac{5}{8}|J_3[(kR_0-i\mu)\theta]|^2\right]. \quad (24)$$

III. COMPARISON OF CALCULATED AND EXPERIMENTAL DIFFERENTIAL CROSS SECTIONS

The formulation presented in the preceding section is now applied to analyze the experimental data on the inelastic scattering of 800 MeV/c charged pions from the 2^+ (4.44 MeV) and the 3^- (9.64 MeV) states of ^{12}C obtained at BNL [2]. The parameters that enter into our calculations are R_0, a, μ, and C_λ defined in Sec. II. The values of R_0, a, and μ have been already estimated in Ref. [8] from the analysis of experimental data on elastic scattering of 800 MeV/c charged pions from ^{12}C obtained at BNL [2]. In order to complete the calculation of the inelastic scattering differential cross sections in which the nucleus is excited from the ground state to one phonon vibrational states we also require the value of the restoring force parameters C_λ. From Eq. (22) we see that C_λ is analogous to a scaling factor that has no effect on the shape of the angular distribution of the cross sections. The values of C_2 for the 2^+ (4.44 MeV) state and C_3 for the 3^- (9.64 MeV) state of ^{12}C have been estimated from their corresponding experimental mean lifetimes which we deduce from the published [15] values of the radiative widths Γ_λ (eV) of 1.08×10^{-2} eV for 2^+ (4.44 MeV) and 3.1×10^{-4} eV for 3^- (9.64 MeV).

To obtain the best fit to the inelastic scattering data, we have varied the values of R_0, a, μ, and C_λ within

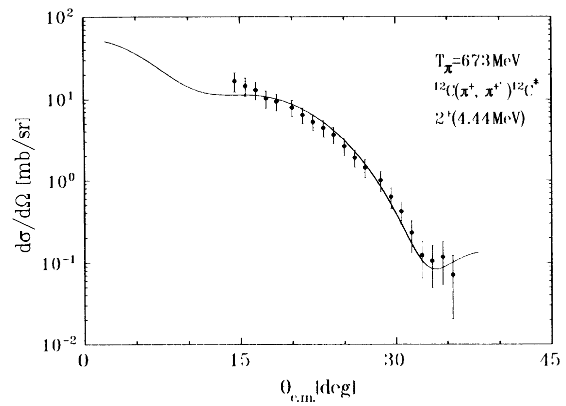

FIG. 1. Comparison of experimental and theoretical differential cross sections of inelastic scattering of positive pions at $p_{\text{lab}} = 800$ MeV/c from the 2^+ (4.44 MeV) state of ^{12}C. The experimental points are taken from Ref. [2]. The result of the present calculation is shown by the solid line.

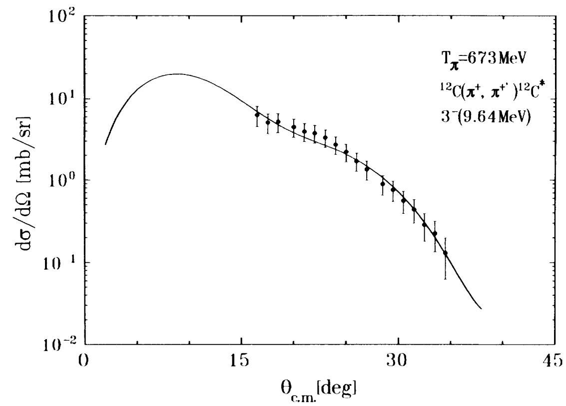

FIG. 3. Comparison of experimental and theoretical differential cross sections of inelastic scattering of positive pions at $p_{\text{lab}} = 800$ MeV/c from the 3^- (9.64 MeV) state of ^{12}C. The experimental points are taken from Ref. [2]. The result of the present calculation is shown by the solid line.

reasonable limits from those values estimated as outlined above. The results of our calculations compared to the experimental results are shown in Figs. 1–4. As can be seen from these figures, the agreement between the calculated and experimental cross sections is in general very good. The values of R_0, a, μ, and C_λ used in our calculation are summarized in Table I. At this point it is relevant to point out that the values of C_λ estimated from experimental mean lifetimes are increasingly sensitive to values of R_0 for states of higher multipolarity as is easily seen from the relation [14].

$$B(E\lambda) = \left[\frac{3}{4\pi} ZeR_0^\lambda\right]^2 \frac{\hbar\omega_\lambda}{2C_\lambda}. \tag{25}$$

For the case of the 2^+ (4.44 MeV) level we see from Table I that the values for R_0, a, and μ leading to the best fit to the experimental data for the π^+ scattering are the same as those estimated from the elastic scattering results. Further we find that the value of C_2 that gives the best fit for the π^+ scattering is the same as that value estimated from the experimental mean lifetime of this state. For the case of π^- scattering for the 2^+ (4.44 MeV) state we again have the same values of R_0 and a that were estimated from the elastic scattering results along with a slightly higher value of μ which is closer to the value obtained for the π^+ case. The estimated value of C_2 has also been used for this case. These results provide very strong confirmation for the underlying vibrational nature of the 2^+ (4.44 MeV) state of the ^{12}C nucleus. In the case of the 3^- (9.64 MeV) state we see that the values of the effective radius R_0 and effective surface thickness a leading to the best fit for both the π^+ and the π^- cases are about 10% higher than our estimates from the elastic scattering. In addition, the value of C_3 utilized to obtain the best fit to the experimental data estimated from the experimental lifetime of the 3^- (9.64 MeV) state corresponds to the nuclear radius which is slightly higher than the value of the radius which leads to the best fit as can be seen from Table I.

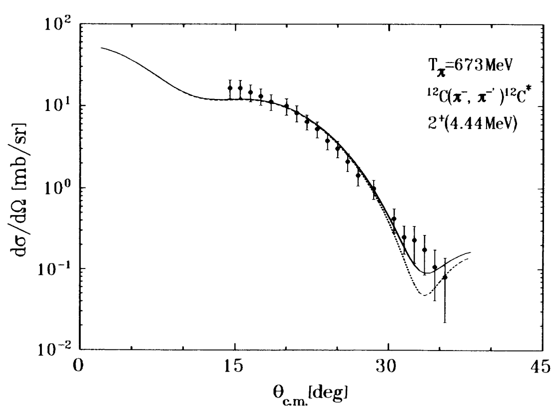

FIG. 2. Comparison of experimental and theoretical differential cross sections of inelastic scattering of negative pions at $p_{\text{lab}} = 800$ MeV/c from the 2^+ (4.44 MeV) state of ^{12}C. The experimental points are taken from Ref. [2]. The solid and dotted curves are the results of the present calculation obtained by using the values of $\mu = 0.55$ and $\mu = 0.35$, respectively.

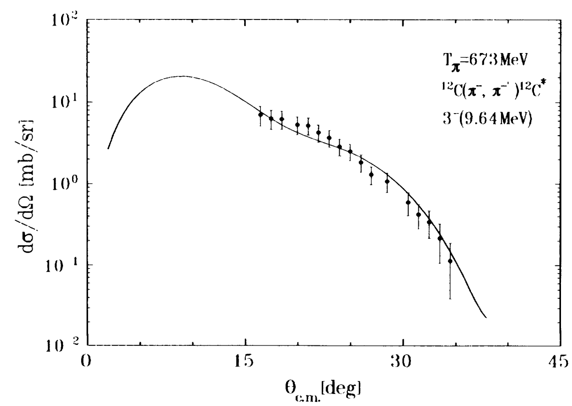

FIG. 4. Comparison of experimental and theoretical differential cross sections of inelastic scattering of negative pions at $p_{\text{lab}} = 800$ MeV/c from the 3^- (9.64 MeV) state of ^{12}C. The experimental points are taken from Ref. [2]. The result of the present calculation is shown by the solid line.

TABLE I. The experimental constants and the parameter values used in the analysis of the elastic and inelastic scattering of charged pions from ^{12}C at $p_{\text{lab}} = 800$ MeV/c.

Probe	^{12}C Level	$k^{\text{c.m.}}$ (fm)$^{-1}$	R_0 (fm)	a (fm)	μ	C_λ (MeV)
π^+	0^+(g.s.)	3.7802	2.38[a]	0.58[a]	0.60[a]	
π^+	2^+(4.44 MeV)	3.7802	2.38	0.58	0.60	18.8[c]
π^+	3^-(9.64 MeV)	3.7802	2.61	0.64	0.60	71[d]
π^-	0^+(g.s.)	3.7802	2.38[b]	0.56[b]	0.35[b]	
π^-	2^+(4.44 MeV)	3.7802	2.38	0.56	0.55	18.8[c]
π^-	3^-(9.64 MeV)	3.7802	2.61	0.61	0.35	71[d]

[a]These values of R_0, a, and μ are obtained from the analysis of the elastic scattering of π^+ from ^{12}C at $p_{\text{lab}} = 800$ MeV/c and given in Ref. [8].
[b]These values of R_0, a, and μ are obtained from the analysis of the elastic scattering of π^- from ^{12}C at $p_{\text{lab}} = 800$ MeV/c and given in Ref. [8].
[c]This value of C_2 is obtained from the experimental mean lifetime of the 2^+ (4.44 MeV) state of ^{12}C using $R_0 = 2.38$ fm.
[d]This values of C_3 is obtained from the experimental mean lifetime of the 3^- (9.64 MeV) state of ^{12}C using $R_0 = 3.02$ fm.

Therefore it is most likely that this state is much more distorted than expected and is not purely collective in character. Indeed it is not surprising that the light ^{12}C nucleus would deviate from purely collective behavior as its angular momentum increases. Evidence for a much less robust collective structure for the ^{12}C nucleus was also seen in the elastic scattering results [8] where the ratio a/R_0 for ^{12}C was seen to be about 0.27 indicating a much more diffuse structure as compared to the much heavier ^{40}Ca nucleus, which was also examined, for which $a/R_0 \approx 0.16$. It would therefore be of considerable interest if more experimental data at a comparable energy range to that of the present study for pions became available for inelastic scattering from heavier nuclei.

IV. CONCLUDING REMARKS

It is clear from the preceding section that the calculations using the phenomenological model in the adiabatic approximation give a very good description of the experimental data from the inelastic scattering of 800 MeV/c charged pions from the 2^+ (4.44 MeV) and 3^- (9.64 MeV) states of the ^{12}C nucleus. The quality of the agreement between the calculated and experimental differential inelastic scattering cross sections shown in Figs. 1–4 strongly supports our assumptions regarding the validity of the phenomenological model calculations. The success of the present analysis as well as that of earlier analyses [8,10,11] tend to suggest that it is mostly the geometrical structure of the target nuclei which determines the elastic and inelastic high energy nuclear scattering. It should further be noted that we have assumed that the 2^+ (4.44 MeV) and 3^- (9.64 MeV) states of the ^{12}C nucleus have the structure of simple one phonon vibrational states. We believe that the results of the present study strongly support a vibrational picture of the 2^+ (4.44 MeV) state. However, the results for the 3^- (9.64 MeV) state indicate that the structure of this state may be much more complex. The results obtained here regarding this state are very much in line with similar results obtained in the earlier study [11] involving 180 MeV antiprotons. However, we believe that the detailed structure of this state does not affect the final results and conclusions obtained within the framework of the present model.

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