THE PHYSICS OF UNPHYSICAL PIONS

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Abstract

This article is a tribute to Hans Weidenmüller. Although our interests have drifted apart over the years, I still fondly remember the time when we both worked on chaotic phenomena in nuclei. Hans always wanted to see the general lesson for other areas of physics drawn out of the special case in nuclear physics, an attitude we both share. Here is an illustration that nuclear physics still has such lessons to offer. It comes from scattering theory to which Hans has contributed so much. The special application is to 0° charge-exchange reactions and their relation to pion physics as well as to neutrino and muon physics.

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1 Introduction

During the last decade a series of nuclear investigations have been based on the use of 0° charge-exchange reactions. The prototype of such reactions in nuclear physics is the (n,p) charge exchange, but by now it is well known that other such reactions at medium energies such as ($^3\text{He},\text{T}$), (d,2p), ($^{12}\text{C},^{12}\text{B}$), or ($^{12}\text{C},^{12}\text{N}$), etc. in most respects mimic the same features and reflect the same physics [1–10]. The nuclear modes that are selected by these reactions are all dominated by pionic features, i.e. by the excitation of nuclear spin-isospin modes for low-energy transfers to the final nucleus (Gamow–Teller modes) and by the predicted pion-like mode for energy transfers characteristic of the Δ resonance [11]. Recently, it has even been reported that single pions have been observed in such reactions and that their forward peaking is so strong that they must have been produced coherently on the target nucleus [12]. These special features are all characteristic of the near-forward direction and they are rapidly lost with increasing transverse-momentum transfer. Such reactions explore the nuclear response for the energy–momentum transfer imposed by the particular kinematics. They are closely linked to the nuclear response to an axial spin-isospin probe and thus to pion physics. The reactions have been extensively analysed in terms of particular models. It is the purpose of the present article to take a look at the physics in a more general perspective. The reason is that the reactions can all be considered to be produced by a ‘beam’ of virtual pions, i.e., pions for which the energy–momentum relation differs from that for free pions $\omega^2 - q^2 = m_\pi^2$. The role of the projectile is simply that of a source of the virtual pion field. It acts only as the carrier of the virtual pion beam, which cannot exist in its absence. Looked at in this way, we are naturally led to consider the characteristic features of a reaction induced by such a virtual beam. This can then serve as a convenient source of inspiration and framework for understanding qualitative features and characteristic differences with reactions induced by free pions. All of the above interactions involve nuclei and it is thus inevitable that the analysis must face the problem of distortion and absorption of the charge-exchange process. It is most interesting that there is one group of closely related reactions for which this problem does not exist, but, alas, it is replaced by formidable experimental difficulties. Neutrino charge-exchange processes at 0° (ν, e^- or μ^-) on nuclei are, by Adler’s theorem [13], exactly related to the scattering and reaction amplitudes induced by a beam of virtual pions with the energy and momentum of the energy–momentum transfer in the process. This is the reason why it has been proposed to investigate the response functions of neutrino reactions [14]. While it is difficult at the moment to explore such neutrino reactions except for very special cases, the undisputed relation to the off-mass-shell pionic amplitudes provides a convenient starting point for the discussion of the physics. A similar situation is also met in the capture of muons from the 1s orbit of muonic atoms. This process corresponds once more to the absorption of virtual pions from a well-understood source and can be discussed accordingly [15]. The following discussion will deliberately be made on a qualitative level, since the purpose is to give a physical insight that can also be used as a guideline both for further problems to be explored and for optimizing experimental conditions.

2 The nature of the virtual-pion beam

2.1 Kinematics

Let us consider one of the projectiles discussed in the charge-exchange reactions above. For the sake of concreteness we take the nucleon. The projectile surrounds itself with a cloud of virtual pions. Let us consider this cloud in the so-called Brick-wall or

Breit frame for the nucleon in which the pion is emitted at no cost in energy but with a momentum k . In the present case this frame nearly coincides with the projectile rest frame. The incident and outgoing nucleons have momentum $\pm k/2$, respectively and the same energy $E_0 = \sqrt{M^2 + (k/2)^2}$. We are concerned only with the case where the pion is emitted along the z -axis, since the transverse-momentum transfer vanishes for 0° reactions. The rest frame energy-momentum transfer is $\omega = 0, q = k$. Let us now boost the projectile into the laboratory frame by giving it a velocity β in the z -direction. The energy-momentum loss in the charge-exchange reaction ω, q is now

$$\omega = \Delta E = \gamma\beta k; \quad q = \gamma k, \quad (1)$$

where $\gamma = \sqrt{1 - \beta^2}$ and the velocity $\beta = (p + p')/(E + E')$. The value for k follows from the relation

$$k^2 = q^2 - \omega^2. \quad (2)$$

In the present case the interesting point is that k is very moderate. For energy losses corresponding to the region of the Δ isobar excitation k is of the order of 240 MeV/ c for a 600-MeV incident nucleon, and the nucleon kinetic energy in the Breit frame is nearly negligible: about 5 MeV. Clearly, by this transformation we have shown the reason why pion effects may be particularly important in such processes: the momentum transfers in the Breit frame are typically on the pion scale even at rather low incident nucleon energies, and the charge exchange also favours pionic transfers. One notes that with increasing energy the necessary momentum k in the Breit frame becomes increasingly smaller for constant energy transfer, since $k^2 = \omega^2(1 - \beta^{-2}) \rightarrow 0$. It is therefore a generic feature of such charge-exchange reactions that any reduction in the cross-sections due to the projectile form factor diminishes as the projectile energy increases.

2.2 The pion source function

The source function for the pion field will of course vary with the projectile as a function of the pion-projectile vertex. Let us take the nucleon as an example. For momentum transfer k it has the form

$$J_\pi(k) = \frac{f}{m_\pi} \sigma_z k \frac{F(k^2)}{F(-m_\pi^2)} \tau^{(+,-)}. \quad (3)$$

Here f is the pseudovector πN coupling constant, $F(k^2)$ the vertex form factor and τ the isospin operator. In the Breit frame this form is identical for pseudoscalar and pseudovector couplings, since there is no energy loss [16]. The explicit appearance of the spin instead of γ matrices is no relativistic restriction. The main difference with the standard coupling to the nucleon or to the $(^3\text{He}, T)$ is the explicit appearance of the form factor with a value for the Breit momentum transfer k which varies as a function of the incident projectile energy and with the energy transfer.

3 Propagation in a uniform medium

Up to this point we have shown that at least part of the interaction in charge-exchange reactions originates in the pion cloud and that the energy-momentum transfer in the reaction singles out the contribution from pions with energy-momentum ω, q such that $\omega^2 - q^2 = -k^2 \neq m_\pi^2$, i.e. there is an incident plane wave pion beam, but the pion

is off the mass shell. Let us now explore the special physical features of such a beam. Its characteristic feature is that it cannot exist unless it is constantly replenished by a source. In order to have an incident wave $\phi_{\text{inc}}(z, \omega) = \exp(iz)$ the free pion wave equation must have a source term:

$$(\nabla^2 - m_\pi^2 + \omega^2)\phi(z, \omega) = (-q^2 - m_\pi^2 + \omega^2)\exp(iz) \quad (4)$$

with $(q^2 + m_\pi^2 - \omega^2) = (k^2 + m_\pi^2)$. As usual, the right-hand side of the equation vanishes when the pion is a physical one. This simply means that a physical pion does not need to be sustained by a source in order to propagate like a free wave. Let us next consider what happens in an infinite nuclear medium. We can then describe the propagation of a coherent pion wave in the medium in terms of a uniform optical potential $U(q^2)$ dependent on the incident wave number q . Since we know that various processes attenuate the coherent wave, we will assume the potential to have an absorptive part, which corresponds to the physics. We then solve the inhomogeneous wave equation

$$(\nabla^2 - m_\pi^2 + \omega^2 + 2\omega U(q))\phi(z, \omega) = (-q^2 - m_\pi^2 + \omega^2)\exp(iz), \quad (5)$$

which has the solution

$$\phi(z, \omega) = \frac{q^2 + m_\pi^2 - \omega^2}{q^2 + m_\pi^2 - \omega^2 - 2\omega U(q)} \exp(iz). \quad (6)$$

Since there is absorption there will be no freely propagating waves in the medium. In general, such waves would be present whenever the dispersion equation has solutions corresponding to a real wave number K in the medium. There is now a wave with the incident wave number q generated throughout the medium. This wave exists for virtual pions only. For real pions it is suppressed by the factor $(q^2 + m_\pi^2 - \omega^2) = 0$, which corresponds to the Ewald extinction theorem for the incident plane wave within a medium. Because of this generated wave there is now an absorption probability per unit volume P

$$P = 2 \operatorname{Im} U(q) \left| \frac{q^2 + m_\pi^2 - \omega^2}{q^2 + m_\pi^2 - \omega^2 - 2\omega U(q)} \right|^2. \quad (7)$$

This means that knowing the normalization of the incident wave, we also know the reaction rate in the medium. Since the absorption processes of the coherent wave of virtual pions is mainly due to either quasi-elastic scattering on nucleons or to real absorption, which is a rather short-ranged phenomenon for physical pions, we have good reasons to expect that the reaction products will be very similar to those for physical pions as long as the energy transfer stays comparable.

Let us next investigate the generation of real pions that occurs in the interface region between the medium and free space. Consider the virtual pion wave propagating through an infinite half-plane to the left, emerging into free space to the right. We recall once more that the medium is an absorptive one. We must now take into account the reflection of the wave at the surface, on the one hand, and the fact that real pions will appear in the free region to the right. Matching the boundary conditions for the problem (to keep notation short we will not discuss the case of a momentum-dependent potential, but this is readily incorporated)

$$\phi_{\text{medium}}(z, \omega) = \frac{q^2 + m_\pi^2 - \omega^2}{q^2 + m_\pi^2 - \omega^2 - 2\omega U(q)} \left(e^{iqz} + \frac{K + q(\omega)}{q + q(\omega)} e^{-iKz} \right). \quad (8)$$

$$\phi_{\text{out}}(z, \omega) = e^{iqz} - \frac{K + q}{K + q(\omega)} e^{iq(\omega)z}, \quad (9)$$

where K is the wave number in the medium and $q(\omega)$ is the vacuum wave number at energy ω of a physical pion. In addition to the wave of virtual pions there is now an additional wave of physical pions with wave number $q(\omega)$. Thus, like the electric transition radiation generated when a charged particle traverses the boundary between two media of different refractive indices, a coherent pion transition radiation is generated in the present case and the physical reason is basically the same one. The missing momentum is furnished by the boundary, i.e. it is furnished by the nucleus and not by the pion momentum distribution, such that it does not reflect in the energy-momentum loss of the projectile carrying the virtual-pion beam. Since the virtual beam is physically distinguished from generated beam, we note that this beam has the amplitude $-(K + q)/(K + q(\omega))$. This means that in the limit of a physical pion the amplitude is -1 , i.e. there is complete shadowing. The emerging beam in the present case will instead correspond to an emission with the the same amplitude but the diffractive effects are closely related: according to Babinet's principle an absorbing disc and an emitting disc produce the same diffraction pattern. We also note from the expression above that the amplitude for emission varies rather weakly with the off-shell nature of the scattering.

3.1 Effects of a pionic mode

Suppose now that there is a mode for the pion in the medium corresponding to the solution of the pion dispersion relation, i.e. that the energy-momentum transfer corresponds to that of a pion propagating spontaneously in the medium. Such a collective pion mode in the medium was already predicted some time ago [11]; it has been clearly realized that it is unsuitable to investigate such a mode using physical pions but that it would show up in investigations of response functions using probes such as neutrinos, charge-exchange reactions etc. We may visualize this mode as the solution of the momentum-dependent optical potential dispersion relation, but the concept is far wider. The wave number is in general a complex one $K = K_r + iK_i$, since the wave will be attenuated. From Eq. (7) it follows that the unphysical pion will have the interaction probability enhanced by a factor $|q^2 - q(\omega)^2 / q^2 - K^2|^2$. This pionic mode will show up as a peak in the energy loss spectrum of the projectile that generates the virtual pion beam. Such a peak has indeed been observed [5], [6] and been interpreted in this fashion [17–20]. It is clear in the approach we have used here that the wave function is generated within the medium with an enhancement factor such that the unphysical pion in general will be absorbed or scattered. This means that not only the energy loss spectrum, but also the spectrum of those reaction products that escape out of the nucleus should be sensitive to the collective mode and show a corresponding enhancement. We may then ask whether the coherently produced pions will also will be enhanced by the mode. The answer to this question is negative. We have already seen that physical pions in scattering processes do not probe the mode due to the mismatch in energy-momentum. For coherently produced pions we see above in our simple model that the amplitude is $-(K + q)/(q(\omega) + q)$. There is no special effect as the value of q satisfies the condition for the pionic mode. This is just another way of saying that a physical pion experiences a mismatch even if it induces reactions. The coherently produced pions are of considerable interest in their own right, but

not in connection with the issue of collective pionic states. Finally, since the projectile is a strongly interacting particle the incident projectile wave is distorted and attenuated on the passage through the nuclear medium quite apart from the pionic interactions. Since the nucleon projectile acts as the source for the virtual pion beam it will also attenuate it. In view of the rather high energy and the fact that we are concerned with scattering at small angles, the Glauber approximation with a straight line trajectory provides a reasonable qualitative guide. This means that for a given impact parameter $\mathbf{b}$ the source is attenuated by a factor $\exp(-\int_{-\infty}^{\infty} \sigma \rho(z, \mathbf{b}) dz)$, where σ is the total NN cross-section. The wave generated in the medium by the virtual pion beam is then attenuated by the corresponding factor.

4 Neutrino reactions

The physics of a virtual pion beam stands out particularly clearly in the case of forward neutrino reactions. Consider a neutrino-induced reaction

$$\nu + N \rightarrow l + X, \quad (10)$$

where N is a nuclear or a hadronic target, l is a lepton (e or μ) and X is the specific final reaction products in the reaction, which can be anything from individual nuclear states, complicated emissions of nucleons, or states with pions whether coherently or incoherently produced. Any one of these specific processes can be related to the corresponding reactions

$$\pi(q, \omega) + N \rightarrow X \quad (11)$$

induced by a virtual plane pion wave carrying the same energy and momentum as the energy-momentum transfer. More precisely: Adler's theorem [13] states that the weak interaction matrix element M for such a forward-neutrino reaction is exactly proportional to the corresponding T -matrix element for the virtual pion-induced reaction.

$$M \propto \frac{G_W}{q_0} T(\pi(q) + N \rightarrow X), \quad (12)$$

where $G_W \simeq 10^{-5} M^{-2}$ is the usual weak interaction coupling constant. Only weak and general assumptions go into this theorem. It assumes on the one hand that the divergence of the axial current is proportional to the pion field (PCAC)[21], which is natural in chiral models in nuclear physics as well, and on the other hand that the lepton mass is negligible, an approximation that is increasingly accurate with increasing neutrino energy. It is therefore evident that 0° is exactly equivalent to the physics of a virtual pion beam with $\omega = q$. In this sense these neutrino reactions are the best possible example of how such virtual beams arise: the neutrino can be viewed as virtually decomposing into a pion and lepton and it is this pion that is responsible for the virtual pion beam. The trouble with the neutrino reactions is the very small cross-sections on the one hand and the poor resolution on the other, which currently make the neutrino a very exotic tool for exploring the physics of virtual pions. This situation may change in the future with the advent of kaon factories. There are still several aspects of the situation for neutrinos that merit further reflection. First, there is the paradoxical situation that neutrino reactions can be proportional to the physics of a strongly interacting particle, which in the physical region has a cross-section proportional to $A^{2/3}$, while neutrinos freely penetrate even a heavy nucleus and have interactions proportional to the nuclear volume or A . The explanation was given by Bell [22],[23]. As we have seen above, the virtual beam generates a non-vanishing wave

throughout a large uniform medium and this wave produces pionic processes proportional to the volume. This means that there is no conflict between weakness of neutrino interactions and their pionic nature. The second interesting statement of the theorem is that the neutrino processes at small energy transfers are also described explicitly by a pion off-the-mass-shell amplitude and that therefore those reactions are pure pion physics. For small space-momentum transfers in such reactions the allowed impulse approximation is a good one. The axial transition interaction is

$$\langle 0 | M^{(+,-)} | n \rangle \propto \sigma_{\text{lepton}} \cdot \langle 0 | \mathbf{A}^{(+,-)} | n \rangle \simeq \sigma_{\text{lepton}} \cdot g_A \langle 0 | \sigma_i \tau_i^{(+,-)} | n \rangle, \quad (13)$$

i.e. the nuclear transition is given by the unretarded Gamow-Teller operator. It is of course well known that the axial transition operator in nuclei is intimately linked to the nuclear pion source function via the hypothesis of axial locality [15], but here we see this statement realized very transparently in a special case.

5 Nuclear muon capture and virtual pion physics

While the Adler theorem and its physics apply to high-energy neutrino reactions, a somewhat similar situation is associated with the physics of the nuclear capture of stopped muons from the 1s orbit of a muonic atom. This is the muonic counterpart of the K-capture of an electron, although for muons the energy available is 106 MeV. The muonic wave function is determined by the extended Coulomb potential. In practice this means that it is nearly constant over the nucleus, which is the only point of relevance for the present discussion. The muon lifetime in this state is long and it is in a heavy nucleus mostly determined by the absorption by the weak-interaction processes. These are dominated by capture due to the axial current coupling, strongly related to pion physics and the excitation of nuclear Gamow-Teller states as has been known for a long time. The muon, being a weakly interacting particle, penetrates freely throughout the nucleus and its wave function is not distorted by strong interactions. However, the muon surrounds itself for a small fraction of the time by a cloud of virtual pions for which it acts as a source. We can view this approximately as produced by the virtual decay of the muon $\mu^- \rightarrow \nu \pi^-$, where the pion has an energy ω corresponding to the energy transfer to the final nucleus. Since the pion wave function is now generated by the constant muon wave function as a source without any attenuation, we have here a particularly clear case of off-the-mass-shell physics. The situation resembles the previous one for neutrinos, but there are some important differences. In order to be precise, we recall that the muon capture per unit energy in the limit of vanishing neutrino energy is proportional to the imaginary part of the off-mass-shell π -nuclear scattering length $a_s(\omega)$ and scattering volume $a_p(\omega)$ [15],[24]. The s-wave capture is associated with the time component of the axial current, the p-wave capture with its space component.

$$\frac{d\Gamma}{dE} \propto \text{Im } a_s(\omega) \omega^{-2} \frac{m_\pi^4}{(m_\pi^2 - \omega^2)^2} + 9 \text{Im } a_p(\omega), \quad (14)$$

where $\omega = m_\mu$. At this particular kinematical point the free-pion wave number is imaginary with $\kappa = (m_\pi^2 - \omega^2)^{\frac{1}{2}}$ and $\omega = m_\mu$ and there is no external space-momentum transfer. As before, the incident wave generated in the off-mass-shell scattering is normalized as $\exp(i\mathbf{q} \cdot \mathbf{r})$, i.e. to unity in the present case. This means that the constant source has been taken as $(m_\pi^2 - \omega^2)$. Consider first the case of a point nucleus. The wave function is then

given in terms of the on-the-mass-shell scattering length:

$$\phi = 1 + \frac{a_s}{1 + \kappa a_s} \frac{e^{-\kappa r}}{r}. \quad (15)$$

In the near-zone of small r this expression reduces to the wave function for a physical pion at threshold in the same situation

$$\phi \rightarrow (1 + \kappa a_s)^{-1} \left(1 + \frac{a_s}{r}\right). \quad (16)$$

Therefore in this case the value for the off-shell s-wave scattering length is

$$a_s(\omega) = \frac{a_s}{1 + \kappa a_s}, \quad (17)$$

$$\text{Im } a_s(\omega) = \frac{\text{Im } a_s}{|1 + \kappa a_s|^2}, \quad (18)$$

which is a well-known result, when the finite size of the nucleus can be neglected. Let us now see what happens for a large nucleus, the interior of which can be treated as an infinite nuclear medium. Although it is not really essential, I will for specificity and transparency of argument assume that the π -nucleus interaction is sufficiently well described in terms of a standard low-energy optical potential of the schematic type:

$$2\omega U = -\nabla \cdot \chi \nabla + Q \quad (19)$$

with the corresponding wave equation

$$\nabla \cdot (1 - \chi) \nabla \phi + Q \phi + (\omega^2 - m_\pi^2) \phi = \omega^2 - m_\pi^2. \quad (20)$$

The basis for this potential is that the interaction in the medium, whether for elastic scattering or for absorptive processes, are local s and p wave processes and the absorption is described by the imaginary part of χ and Q . As long as the unphysical pion has an energy ω not too far below the physical pion mass and for momenta that are not large, this description is a good approximation. We recall that the local interaction Q is repulsive, while the velocity-dependent one in χ is attractive. In nuclear matter the generated wave function is a constant. Since $\nabla(\text{constant}) = 0$, the generated wave is

$$\phi_0 = \frac{\omega^2 - m_\pi^2}{\omega^2 - m_\pi^2 + Q}. \quad (21)$$

The remarkable result is therefore that the momentum-dependent terms are irrelevant and that in this case there is *no* local p-wave interaction in the medium. The local absorption rate in the medium per unit volume is

$$P = 2 \text{Im } U(\mathbf{q}^2 = 0) | \phi_0 |^2 = \frac{\text{Im } Q}{\omega} \left| \frac{\omega^2 - m_\pi^2}{\omega^2 - m_\pi^2 + Q} \right|^2. \quad (22)$$

From this we can write the absorptive part of the scattering length at an energy ω in terms of the integral of the constant rate per unit volume P over the volume V :

$$\text{Im } a_s = V \frac{\text{Im } Q}{4\pi} \left| \frac{\omega^2 - m_\pi^2}{\omega^2 - m_\pi^2 + Q} \right|^2. \quad (23)$$

The corresponding discussion for the p-wave scattering volume a_p is immediate and follows nearly identical reasoning. For a small nucleus we have

$$a_p(\omega) = \frac{a_p}{1 + \kappa^3 a_p}. \quad (24)$$

For an extended medium we use a small value for the unphysical momentum $\mathbf{q}$ to single out the p wave. This means in the limit of vanishing momentum that the imaginary scattering is given by a volume integral over $\text{Im } \chi$ of the same type as before:

$$3 \text{ Im } a_p = V \frac{\text{Im } \chi}{4\pi} \left| \frac{\omega^2 - m_\pi^2}{\omega^2 - m_\pi^2 + Q} \right|^2. \quad (25)$$

From these expressions we see at once that the Born approximation holds when the interaction Q is small compared to the characteristic squared wave number $m_\pi^2 - \omega^2$. The reason is that then the external driving force imposes itself completely. The condition for the constant wave in the medium is that the characteristic distance of the wave number in the medium should be smaller than the radius, i.e. that $\sqrt{(m_\pi^2 - \omega^2 - Q)/(1 - \chi)} R < 1$. This means that the range of the Yukawa field about a point source in the medium should stay inside the medium. The pointlike approximation requires two conditions to be valid. First, κR must be smaller than unity, which simply means that the free Yukawa factor must have a range larger than the nucleus. In addition there is the hidden assumption that the interaction potential is larger than the $\pi\mu$ mass difference. We can now insert these expressions (22) and (24) for a large nucleus into the expression for the end-point capture rate (14). This gives

$$\frac{d\Gamma}{dE} \propto \text{Im } Q \omega^{-2} \frac{m_\pi^4}{|m_\pi^2 - \omega^2 - Q|^2} + 3 \text{ Im } \chi \left| \frac{\omega^2 - m_\pi^2}{\omega^2 - m_\pi^2 + Q} \right|^2. \quad (26)$$

We therefore see clearly, in the situation where the energy transfer to the nucleus is maximal, that the problem of the muon capture in a large nucleus is exactly described by the strong-interaction physics of off-mass-shell pions. What about μ -capture into excited nuclear states in the first 10–20 MeV above the ground state? This is the region in which the bulk of the capture takes place and the momentum transfer by the neutrino is typically of the order of 100 MeV. The reaction is dominated by the axial-current capture. The description of these processes is in a sense ‘trivial’ in so far as an impulse approximation on the individual nucleons is an excellent first approximation. The axial current is therefore described by the standard spin–isospin operator with momentum transfer and the predominant excitations are the spin–isospin analogues of the giant dipole resonance. On the other hand, as for neutrino reactions, the axial current can also here be linked directly to the pion source function. The impulse approximation is just the approximation to virtual-pion physics of very low energy ω in which the individual nucleons have the usual $\sigma \cdot \nabla$ coupling. The more sophisticated description in terms of axial locality to this situation generates naturally higher-order many-body terms to the nuclear axial current due to pion propagation in the nuclear environment [15],[25]. These are the terms usually referred to as meson exchange currents, which become predominant in situations of large energy transfer.

6 Conclusion

It is a generic feature of a number of forward-charge exchange reactions that their physics can be viewed as produced by a beam of virtual pions. The characteristic feature

of such ‘unphysical’ pions is that a pion wave can be generated by the projectile also inside the nuclear medium. This is a region that is poorly accessible to physical pions for large nuclei. In particular, physical pions are not suitable for the exploration of the collective pion branch, the pionic mode, because of the mismatch in the pion dispersion relation. The virtual pions, on the contrary, can display such a mode both via the energy-loss spectrum of the projectile and via the yield of reaction products which will be enhanced when the energy–momentum condition of the pionic dispersion equation in the medium is fulfilled. Coherently produced physical pions are a natural feature of such virtual pion beams and they are generated by a mechanism similar to the one that gives rise to the transition radiation which occurs when a charged particle traverses the boundary between two media of different refractive index. In the same spirit, it is unnatural to treat the nuclear spin–isospin states excited by forward-charge exchange in a different perspective from the higher excitations in the Δ region with the pion collective state and pion production as well as the associated disintegration phenomena. It is a common procedure to ascribe the dominant excitation of Gamow–Teller states in forward-charge exchange reactions to the spin–isospin part of the nucleon–nucleon interaction. The mechanism is assumed to be a distorted wave impulse approximation for which the contributions from the struck nucleons in the target are added to give an effective Gamow–Teller transition operator. While this reasoning can be defended, it misses most of the beauty and generality of the situation. It is well known that the central spin–isospin interaction in np charge exchange is intimately related to the one-pion-exchange (OPE) interaction. The natural way of viewing these exchange reactions, whether for nucleons, neutrinos or heavy ions, is in terms of pion physics. The special approaches in the literature should be viewed as approximations to this general picture. The physics of this situation stands out with special clarity for the neutrino-induced reactions. Consequently, such reactions are of great fundamental interest in spite of the intrinsic experimental difficulties. This is so not only because of the special mechanism linked to the generation of virtual pions but also because of the nuclear information they can potentially yield on the nuclear axial-current response [14].

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