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CHIRAL FILTER, AXIAL CHARGES AND GAMOW-TELLER STRENGTHS

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Abstract - The different ways that nuclear matter responds to the weak axial-vector current are interpreted in terms of modification of the "vacuum" in baryon-rich environments. The notion of "chiral filter" is introduced. Use of a Ward identity is suggested. The Gamow-Teller quenching and the enhanced axial charge in $0^+\rightarrow 0^-$ transitions follow from this. I also discuss briefly possible relevance of the nucleon as a topological soliton configuration to the global property of nuclear axial response functions.

I - INTRODUCTION

How nuclear matter responds to the weak axial current $A_\mu^i(x)$ is a particularly fascinating subject in view of the recent developments in the theory of the structure of the nucleon on the one hand and in experimental observations on the other hand.

Consider the axial vector coupling constant $g_A$ in nuclei. When suitably defined, one can discuss this in terms of an effective axial charge in nuclear medium /1/. I will denote this $g_A^{\text{eff}}$. This quantity has been measured in Gamow-Teller transitions, both in nuclear beta decays /1/ and more recently in giant Gamow-Teller resonances /2/. A surprising observation is that $g_A^{\text{eff}}$ is close to 1. In fact, in a beautiful analysis of the magnetic moment and beta decay data for $3 \leq A \leq 41$, Buck and Perez /3/ show

$$g_A^{\text{eff}} = 1.00 \pm 0.02$$

as contrasted with the free space value

$$g_A = 1.25 \pm 0.01.$$  \hspace{1cm} (2)

The data on giant Gamow-Teller resonances /2,4/ extend this to $A \approx 208$, and confirm that the axial charge is effectively quenched in nuclear matter.

Instead of measuring $g_A^{\text{eff}}$ with the space component of $A_\mu^i(x)$ as in the Gamow-Teller transition, one might consider measuring it with the time component $A_0(x)$. Here the relevant transition would be
Unlike in the vector current, there is a subtle problem of getting at the "charge" with the "charge density" operator $A_0(x)$ (this is the problem faced in deriving $g_A$ from the charge commutator in the way that the Adler-Weisberger relation is derived). So, I have to be more careful in defining it. The single-particle operator (so-called impulse approximation) for $A_0(x)$ is of the form

$$A_0^{S.P.}(x) = \frac{g_A}{M} \sum_i^{+} \gamma_i \cdot 0_i \cdot p_i \delta(x-x_i).$$

Now let me write the full $A_0(x)$ effective in nuclei as

$$A_0(x) = \frac{g_A^{eff'}}{g_A} A_0^{S.P.}(x).$$

[I have put a prime here to distinguish it from (1)].

It has been shown in a more or less model-independent way /5/ that

$$g_A^{eff'} \approx 1.4 - 1.6$$

and the available data /6/, combined with theoretical analyses /7/, seem to corroborate the prediction (6). (I will argue later that the prediction is firm on the basis of chiral symmetry). Thus this measurement indicates that an effective axial charge $g_A^{eff'}$ is considerably enhanced

$$g_A^{eff'} \approx 1.8 - 2.0.$$  

The purpose of this talk is two-fold: first to emphasize that understanding the quenching (1) on the one hand and the enhancement (7) on the other hand which may very well be one of the most fundamental problems in nuclear physics may reveal the complex structure of the strong interaction "vacuum"; and second, to propose a (tentative) mechanism to explain these phenomena, using chiral symmetry considerations /8/.

I will make a brief encounter with our recent views of the nucleon as a topological soliton in connection with the role of the isobar $\Delta$ (1232).

II - NUCLEUS IS A CHIRAL FILTER

Chiral symmetry, a symmetry of QCD associated with (nearly) massless $u$ and $d$ quarks relevant to nuclear physics, is believed to be realized in nature in the Goldstone mode, the triplet of pions being the requisite Goldstone bosons. There are ample evidences for this for nucleons in isolation. But how is the symmetry manifested inside nuclear matter? Recent Monte Carlo calculations /9/ on lattice confirm that when nucleons are jammed into much higher densities than that of nuclear matter, a phase change occurs: chiral symmetry is realized (in high densities) in the manifest Wigner-Weyl mode (or sometimes referred to as quark-plasma phase). The question is: what is the vacuum structure (vis-a-vis chiral symmetry) inside nuclear matter?

Nobody has a clear answer to this. I made a tentative attempt to answer this question at the 1982 Telluride meeting /8/.

Consider the response of a nucleus (in particular, two-body system*) to the vector

* Recent studies /16/ indicate that it suffices to confine to a two body, even for many-body systems.
current \( V_\mu(x) \) or the axial-vector current \( A_\mu(x) \). Using general arguments based on chiral symmetry /10,12/, one can describe systematically the current \(-N-N \rightarrow N-N\) amplitude for small momentum transfer in terms of a soft-pion term and non-soft-pion terms including loop corrections. In diagrams

\[
\begin{align*}
\text{Current} & \quad \pi \quad \text{Soft-\pi Term} \quad + \\
\text{Non-soft-\pi Term} & \\
\end{align*}
\]

Fig. 1

One can write the non-soft-pion terms in an explicit form, given a chiral \( SU(2) \times SU(2) \) Lagrangian, in terms of the parameters of Lagrangian (e.g. \( F_\pi \) = pion decay constant) and quantities depending on \( \mu \), where \( \mu \) is an arbitrary renormalization scale and \( Q \) is an effective momentum carried by the pions.

Now the notion of the "chiral filter" can be stated as follows: whenever the vertex (current + \( N \rightarrow \pi + N \)) is unsuppressed by kinematics and/or symmetries, the soft-pion term is dominant while the non-soft-pion terms are suppressed, whereas whenever it is suppressed by kinematics or symmetries, the non-soft-pion terms are large, with no individual contribution dominating the series.

The specific examples illustrating this notion: the time component of the axial current (namely the axial charge density \( A_0(x) \)) and the space component of the vector current (namely, the nuclear \( M_1 \) operator) are predicted to be dominated by the soft-pion term, whereas the space component of the axial current (namely, the Gamow-Teller operator) and the time component of the vector current (the charge form factor) are expected to be dominated by non-soft-pion terms /5/.

Two cases support (so far) the first point: the electrodisintegration of deuterons at large momentum transfers (\( M_1 \) operator) /12/ and the \( 0^+ \rightarrow 0^- \) (\( \Delta I = 1 \)) beta transitions \( [A_0(x)] /6,7/ \). The recent work by Mathiot and Riska /12/ has cleared up a confusing situation in the interpretation of electrodisintegration process: they show that when consistently treated there is a surprising cancellation between hadronic form factors and heavy-meson exchanges in such a way that the soft pion term dominates up to a momentum transfer \( q > 5 f_{\pi}^{-1} \). The agreement between theory and the experiments remains very good. This work (of Mathiot and Riska) illustrates a simple mechanism for near complete suppression of non-soft-pion terms. The situation is quite similar in the \( 0^+ \rightarrow 0^- (\Delta I = 1) \) transition. The soft-pion term enhances the \( g_{\text{eff}}^A \) by the amount indicated in (6) and (7), this enhancement increasing the decay rate by a factor of about 4. This has been more or less confirmed by refined experiments /6/ combined with a sophisticated calculation /7/. Furthermore a calculation by Kubodera and his coworkers /13/ explicitly verified that the non-soft-pion corrections are insignificant, at most \( \lesssim 10\% \) of the soft-nion term. As a whole, both theoretical and empirical evidences are strong that the loop terms are highly suppressed relative to the soft-pion contribution. This may be interpreted physically as an enhanced participation of the pion cloud in the Goldstone phase.

The Gamow-Teller transition illustrates the second point of the chiral filter notion (I have nothing to say about the charge form factor). Whenever the soft-pion is screened (by kinematics and symmetry), it seems that it is totally screened. The precise way that this happens is yet unknown, though some conjecture will be a subject of later discussions. However it seems clear that this happens whenever
short-range encounter between two nucleons takes place. The Gamow-Teller probe is specific to the short-range spin-isospin correlation $\tau_1 \cdot \tau_2 \cdot \sigma_1 \cdot \sigma_2$ and seems to sample the nuclear environment that is basically different from that of $M_1$ and axial charge-density operators. This spin-isospin correlation is closely related to the topological soliton structure of the nucleon, first discovered by Skyrme /14/ and recently rediscovered by several people /15,16/.

3 - AXIAL WARD IDENTITY AND QUENCHING OF $g_A$.

The phenomenon in question reflects the manifestation of chiral symmetry in nuclei or more specifically in nuclear matter. Although the standard nuclear theory - the shell model - is highly suitable for treating phenomena associated with the Wigner-Weyl symmetry (e.g. vector currents), it is not at all clear that it can correctly describe the Nambu-Goldstone symmetry without some basic additions or modifications.

The approach I suggest is field theoretic /17/. Assume that the axial current is exactly conserved (CAC). (Of course all the low-energy theorems can be derived in CAC; one just has to keep track of the zero-mass pion poles). We will imagine doing a field theory for nuclear matter with a Lagrangian that is then SU(2) x SU(2) symmetric. Let $\Gamma_{\mu}^i(p_f,p_i;\rho)$ be the full unrenormalized vertex for the response to the axial current $A_\mu$ for a suitably defined quasiparticle state $i$ making a transition to a quasiparticle state $f$, with a momentum transfer $p_i - p_f = Q$. Let $S(p;\rho)$ be the full unrenormalized quasiparticle propagator. Then we have the Ward identity, which follows from $\partial^\mu A_\mu = 0$:

$$ (p_f - p_i)^\mu \Gamma_{\mu}^i = S(p_f)^{-1} \frac{1}{2} \gamma_5 \Gamma_{\mu}^i S(p_i)^{-1}. $$

(8)

This is valid for all density $\rho$ including the free nucleon ($\rho = 0$). The theory has the usual ultraviolet divergences, so the usual renormalization has to be made to render the physical quantities finite. In addition, one can make suitable finite renormalizations which one may do by cutting off states beyond some given energy (i.e. truncations). All these have to be done in such a way that the Ward identity (8) is satisfied at each stage. (The renormalization constants must also satisfy the renormalization group property). Let me denote the fully renormalized quantities* by tilde and the density-dependent renormalization constants* by $Z$ as

$$ \tilde{\Gamma}_{\mu}^i(p_f,p_i) = Z_A^{-1} \Gamma_{\mu}^i(p_f,p_i) $$

(9)

$$ S(p) = Z_{S}(p) $$

Defined such that as $\rho \rightarrow m_{\text{eff}}$ ($m_{\text{eff}}$ = effective mass of the nucleon) and $Q_\mu \rightarrow 0$,

$$ (\Gamma_{\mu}^i(p,p) - \text{pion pole term}) + \frac{1}{2} \gamma_5 \mu \gamma_5 S(p)^{-1} \rightarrow \rho - m_{\text{eff}} $$

* Unless otherwise stated, all the quantities written down are density dependent.
Thus in terms of renormalized quantities

$$ (p_f - p_i) \mu \, S_{\mu}^{A \mu} = \frac{Z_A}{Z_2} \left[ \frac{g^2}{\pi^2} \left( \frac{1}{2} \gamma_5 \gamma_\mu + \frac{1}{2} \gamma_\mu \gamma_5 \right) \frac{(p_f + p_i)}{(2m)} \right] $$

Unlike in the vector current case where $Z_v/Z_2 = 1$ follows by letting $Q \rightarrow 0$, $Z_v/Z_2$ is not in general equal to 1 because of the pion pole term in $S_{\mu}^{A \mu}$.

In fact, if one calculates the axial-vector current matrix element in the limit $Q \rightarrow 0$, one finds that

$$ (\frac{Z_2}{Z_A}) = (\frac{g_{\text{eff}}}{g_A}) $$

This is formally analogous to the vector-current where the cancellation is exact for any density. The latter is nothing but a consequence of the charge conservation, as for example the equality of the proton and positron charges, and is seen in all exact symmetries realized in the Wigner-Weyl mode. In the axial current case, the cancellation must occur in a much more subtle way.

It is inconceivable that the nuclear matter is already in the Wigner-Weyl mode of chiral symmetry. On the other hand, from the Ward identity point of view, once the probe is inert to the pion degrees of freedom, there may be no clear distinction between the two modes. Suppose the probe sees two nucleons coalesced into a 6-quark system, though not necessarily in one confinement region (or bag). Then apart from a fine-structure effect, the configurations NN and NA cannot be distinguished. In such a case, a simple consideration based on the Adler-Weisberger relation shows that $g_A^{\text{eff}}$ should be $\sim 1$. This way of describing the phenomenon may not be different physically from saying that far away from the bag surface, a d-quark decays to a u-quark with $g_A = 1$ and the Gamow-Teller operator is blind to confinement, pion cloud etc ...

In practice, $Z_2 = Z_A$ in nuclear matter is a very stringent constraint. Thus there must occur intricate cancellations between possibly large terms. Such cancellations between finite cut-off-dependent terms were already conjectured a long time ago /18/.

* It would be extremely interesting to see how this comes about explicitly, using, say, the linear $\sigma$-model.
without any Ward identity arguments*, and are recently put into a Ward identity result by Zhu and Wong /19,17/. Some cancellations are observed in shell-model calculations /20/, but a (near) complete cancellation cannot be easily assured in standard calculations. It is in my opinion an open problem how to implement chiral invariance in nuclear structure calculations, in particular in the framework of the shell model. It should be recognized that chiral symmetry is intrinsically a relativistic phenomenon.

4 - ROLE OF THE ISOBAR \( \Delta(1232) \) IN THE QUENCHING OF \( g_A \)

The description of the quenching of \( g_A \) in terms of coupling to an "elementary" \( \Delta \) was first introduced in 1974 /21/ and the argument has subsequently been refined /22/ in connection with giant Gamow-Teller resonances*. In terms of Landau-Migdal parameter \( (g'_0)_{\Delta} \equiv g'_{\Delta A} \), the result is

\[
\frac{g_A^{\text{eff}}}{g_A} \approx \frac{1}{1 + \alpha g'_{\Delta A}}
\]

where \( \rho \) is the average density; \( \alpha \) is a constant. If one assumes the "universality relation" /22,23/

\[
(g'_0)_{NN} = (g'_0)_{\Delta A} \equiv g'_0
\]

and for accepted values \( g'_0 = 0.6 - 0.7 \), \( g_A^{\text{eff}} \) so calculated comes into agreement with experiments.

A few comments are in order here:

(1) Some people /24/ have argued that because of the large tensor force-mediated core polarization, the \( \Delta \) plays a minor role in the quenching. It cannot be ruled out that this is still a correct argument, but I believe that it is wrong. In fact, the large core polarization is an artifact of the cut-off-dependent renormalization effect and should be suitably cancelled by other terms through the Ward identity. (An initial attempt to do so is found in Refs. 17 and 19). This criticism applies to other esoteric mechanisms /25/ proposed to account for the quenching.

(2) Gerry Brown has proposed a very simple model for \( g'_0 \) involving \( \pi \) and \( \rho \) exchanges in the presence of short-range correlations /26/. In his model, the universality was automatic. It has recently been argued /24/ that if one takes the \( \pi \) - and \( \rho \) - exchange forces, and calculates \( g'_0 \), because of the exchange term which is more important for the \( \Delta \)-hole channel than for the nucleon-hole channel, the universality relation is destroyed. Gerry will show in his talk that the exchange term in the \( \Delta \)-h channel should actually be strongly suppressed because of the screening ignored in the work of others. I will give another reasoning based on the soliton structure of nucleon which suggests at least at non-perturbative level that the universality relation should not be modified significantly.

At the moment, I see no reason to doubt that the \( \Delta \) plays a predominant role in quenching \( g_A \). This mechanism is consistent with the description based on the Adler-Weisberger relation and the notion of the chiral filter /8/, and also with the recently rediscovered soliton picture of the N and the \( \Delta \).

* Previous arguments based on the Ward identity are found in Ref. 17

**Quenching in terms of a correlation hole (or the Lorentz-Lorenz effect) was discussed by Ericson et al. (Phys. Lett. 47B (1973) 381).
Skyrme's idea /14/ of 1960 that the nucleon and the A arise as a topological soliton in the non-linear σ model has recently been revived /15,16/ in the light of QCD, particularly in the chiral bag model /16/. In this picture, the N and the A are a soliton with a common mass $M_0$ which is inversely proportional to the coupling constant $\alpha = N_c^{-1}$, $N_c$ being the number of color in QCD. They are split off when the soliton is slowly rotated, the splitting $J(J+1)/2\hbar$ being of $O(1/N_c)$.

What is the implication of this description on $g_1^0$? Two aspects are relevant for this quantity. First at the soliton level, namely to $O(N_c)$, the spin and isospin are highly correlated (more specifically $K = J + T$, the diagonal subgroup, is the invariance of the theory) and hence the prediction is consistent with a strong spin-isospin channel (i.e. $g_1^0$). Secondly, the $g_1^0$ represents a short-range correlation and hence is related closely to the repulsive core. To leading order $O(N_c)$, the repulsive core is identical in the N and A.

Therefore at the soliton level, we expect

$$\langle g_1^0 \rangle_{NN} = \langle g_1^0 \rangle_{NA} \cdot$$

This is also consistent with the argument that in terms of 6-quark configurations, if one ignores $O(\alpha_S)$ effect, an NA configuration is as easily excited as an NN configuration.

I do not know how to reliably compute $O(1/N_c)$ corrections to the universality relation; it is conceivable that $O(1/N_c)$ corrections will introduce some deviations. A rough estimate can, however, be made as follows: the $O(1/N_c)$ correction which splits the A and the N gives $\Delta M = M_A - M_N$. From Ref. 27, I estimate

$$\frac{\delta g_1^0}{g_1^0} \sim \frac{\Delta M}{M_0} \sim 0.2 \cdot$$

I expect that this will make the $\langle g_1^0 \rangle_{NA}$ bigger than the $\langle g_1^0 \rangle_{NN}$ by at most about 20%. In the absence of more fundamental calculations, it seems safe to take the same $g_1^0$ to leading order. An experimental verification would give us a valuable clue as to how non-leading terms can be calculated.

6 - AXIAL CHARGE IN NUCLEI

Recall that the "axial charge" $g_A^{\text{eff}}$ measured with the time component of the axial current is different from that ($g_A^{\text{eff}}$) measured with the space component of the axial current; relative to the free nucleon value, the former is enhanced whereas the latter is quenched. One way of understanding this difference is to recognize that although conserved in the world of zero-mass u and d quarks, the axial charge density $A_0(x)$ is not an invariant quantity. The $0^- \leftrightarrow 0^+(\Delta I=1)$ transition samples the axial charge density at rest frame, infested with "soft" phenomena, in particular with soft pions; in this frame, the axial charge would appear enhanced.

Suppose now we boost the system to an infinite-momentum frame at which the axial charge becomes an invariant quantity; in this frame, all the "soft" stuff evaporates, so the effect of the soft pion disappears /28/. Specifically, the soft-m ex-
change current in Fig. 1 goes to zero as \( O(\frac{Q}{p}) \) as \( p \to \infty \) for \( Q \) fixed. In this frame, the axial charge measured with \( A_0(x) \) would appear quenched. Now the Gamow-Teller operator measures the invariant charge. Therefore the quenching seen in Gamow-Teller transitions must therefore be quenched. (Of course the transition rate is invariant under the choice of the frame).

This consideration suggests an intriguing possibility that different nuclear operators sample different facets of the role of chiral symmetry, in particular through the manifestation of the Goldstone bosons \( [ \text{The recent anomaly observed in deep inelastic structure functions of nuclei (the EMC effect) can perhaps be explained in a similar manner}] \). Nuclei thus seem to offer a unique laboratory to explore many facets of the complex QCD vacuum.

I am very grateful for continuous and fruitful discussions with my colleagues, G.E Brown, J. Delorme, A.D. Jackson, F.C. Khanna, K. Kubodera, I.S. Towner and V. Vento.

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