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Gauged Linear σ-Model and Nucleon Structure*

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Abstract

The gauged linear σ-model brings an explicit quark sector to the gauged nonlinear σ-model. Interaction in this sector can lead to a condensed ground state of chiral quarks. The nucleon emerges as a topological soliton surrounded by a chiral condensate. High energy elastic pp and $\bar{p}p$ scattering data from CERN ISR and SPS Collider appear to provide strong evidence in favor of this nucleon structure.

I. INTRODUCTION

Last three decades of intensive research on the structure of the nucleon following the birth of the quark model has not led to a unique description of the nucleon. Instead, it has generated a wide array of models:[1–3] the constituent quark model, the MIT bag model, the chiral bag model, the topological soliton model, the non-topological soliton model, the color dielectric model etc. Experimentally established low energy properties of the nucleon and low energy nucleon-nucleon interaction have not been able to single out one model as incorporating the key features. Surprisingly, from high energy elastic proton-proton and antiproton-proton scattering, there seems to be strong evidence in favor of one of the models: namely, the gauged nonlinear σ- model that describes the nucleon as a topological soliton and introduces the vector mesons ω, ρ, A_1 as gauge bosons.[4] This model, however, needs to be extended (i.e., made linear) to include an explicit quark sector, where left and right quarks can interact via a scalar field and can form a condensed ground state (a chiral condensate).[5] The nucleon appears as a topological soliton surrounded by a chiral condensate.[6,7]

II. GAUGED NONLINEAR σ-MODEL

Our first step is to understand how the gauged nonlinear σ-model is built. To this end, we begin with a much simpler model—the Gell-Mann–Levy σ-model with $SU(2)_L \times SU(2)_R$ global symmetry and Higgs potential. The model is described by the Lagrangian

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$$\mathcal{L} = \bar{\Psi} i\gamma^\mu \partial_\mu \Psi + \frac{1}{2}(\partial_\mu \sigma \partial^\mu \sigma + \partial_\mu \boldsymbol{\pi} \cdot \partial^\mu \boldsymbol{\pi})$$
$$-G\bar{\Psi}(\sigma + i\boldsymbol{\tau} \cdot \boldsymbol{\pi}\gamma^5)\Psi - \lambda(\sigma^2 + \boldsymbol{\pi}^2 - f_\pi^2)^2 \ , \tag{1}$$

where Ψ is the fermion or quark field and f_π is the pion decay constant. We next introduce the unitary field $U(x) = \exp(i\tau^a \phi^a(x)/f_\pi)$, and write $\sigma(x) + i\boldsymbol{\tau} \cdot \boldsymbol{\pi}(x) = \zeta(x)U(x)$. Here, $\phi^a(x)$'s (a= 1,2,3) are the pions or the Goldstone bosons, and $\zeta(x)$ is a scalar field ($\sigma^2 + \boldsymbol{\pi}^2 = \zeta^2$). Eq.(1) can be written as

$$\mathcal{L} = \bar{\Psi}_L i\gamma^\mu \partial_\mu \Psi_L + \bar{\Psi}_R i\gamma^\mu \partial_\mu \Psi_R + \frac{1}{4}\zeta^2 \mathrm{tr}[\partial_\mu U \partial^\mu U^\dagger] + \frac{1}{2}\partial_\mu \zeta \partial^\mu \zeta$$
$$-G\zeta(\bar{\Psi}_L U \Psi_R + \bar{\Psi}_R U^\dagger \Psi_L) - \lambda(\zeta^2 - f_\pi^2)^2 \ , \tag{2}$$

where $\Psi_L = \frac{1}{2}(1-\gamma^5)\Psi$ and $\Psi_R = \frac{1}{2}(1+\gamma^5)\Psi$ are the left and right chiral fields.

The model given by Eq.(2) can be easily generalized to $SU(3)_L \times SU(3)_R$ symmetry with local gauge invariance by considering the unitary field $U(x) = \exp(i\lambda^a \phi^a(x)/f_\pi)$ and by introducing the left gauge field $\mathcal{A}_\mu$ and the right gauge field $\mathcal{B}_\mu$ (λ^a's are the Gell-Mann matrices; a=1–8). We now write a path-integral representation for the partition function of this generalized model:

$$Tr\left[e^{-iHT}\right] = \frac{1}{\mathcal{N}} \int DUD\zeta D\psi D\bar{\psi} \exp\Big[i \int d^4x \Big(\frac{1}{4}\zeta^2 tr[\mathcal{D}_\mu U(\mathcal{D}^\mu U)^\dagger]$$
$$+\frac{1}{2}\partial_\mu \zeta \partial^\mu \zeta - V(\zeta) + \bar{\psi}^{\mathcal{A}} i \not{\mathcal{D}}(\mathcal{A}) a_- \psi^{\mathcal{A}} + \bar{\psi}^{\mathcal{B}} i \not{\mathcal{D}}(\mathcal{B}) a_+ \psi^{\mathcal{B}}$$
$$-G\zeta[\bar{\psi}^{\mathcal{A}} U a_+ \psi^{\mathcal{B}} + \bar{\psi}^{\mathcal{B}} U^\dagger a_- \psi^{\mathcal{A}}]\Big)\Big] \ ; \tag{3}$$

here, $\mathcal{D}_\mu U$ is the covariant derivative of U: $\mathcal{D}_\mu U = \partial_\mu U + \mathcal{A}_\mu U - U\mathcal{B}_\mu$; $V(\zeta) = \lambda(\zeta^2 - f_\pi^2)^2$, $\not{\mathcal{D}}(\mathcal{A}) = \gamma^\mu(\partial_\mu + \mathcal{A}_\mu)$, $\not{\mathcal{D}}(\mathcal{B}) = \gamma^\mu(\partial_\mu + \mathcal{B}_\mu)$, $a_\pm = \frac{1}{2}(1 \pm \gamma^5)$.

The path-integral representation shows how quantum effects come into the picture. The fermion measure in Eq.(3) is gauge dependent and can be written as the product of a Jacobian and an invariant measure:[5]

$$D\psi D\bar{\psi} = \mathcal{J}[\mathcal{A}, \mathcal{B}, U] D\psi^o D\bar{\psi}^o \ . \tag{4}$$

The Jacobian can be identified as $\exp[i\Gamma_{WZ}]$, where $\Gamma_{WZ} = \Gamma_{WZ}[\mathcal{A}, \mathcal{B}, U]$ is the Wess-Zumino action. The net result, therefore, is the appearence of an additional piece of action in Eq.(3). It is not surprising that currents which were conserved before are no longer conserved, because of the contribution of Γ_{WZ}. This is known as the anomaly; namely, a current that is conserved at the classical level is no longer conserved when quantum effects are taken into account.[8]

A simple model can now be built if we: (i) replace the scalar field $\zeta(x)$ by its vacuum value f_π from the very beginning, (ii) keep the WZ action arising from the fermion measure, and (iii) disregard the interaction between the left and right quarks. This is the nonlinear σ-model described by the partition function:

$$Tr\left[e^{-iHT}\right] \simeq \frac{1}{\mathcal{N}'}\int DU \exp\left[i\int d^4x \frac{1}{4}f_\pi^2 tr[\mathcal{D}_\mu U(\mathcal{D}^\mu U)^\dagger] + i\Gamma_{WZ}[\mathcal{A},\mathcal{B},U]\right] . \quad (5)$$

To describe the nucleon using this model, one further includes an additional $U(1)_V$ gauge invariance connected with the baryonic charge conservation, and one introduces the vector meson ω as the gauge boson coupled to the baryonic charge. Limiting oneself to the $SU(2)_L \times SU(2)_R$ sector, one now has ω, ρ, and A_1 introduced as gauge bosons and a WZ action that involves these bosons.[4,9]

The WZ action shows that the ω couples to the pseudoscalar mesons via the interaction $g_\omega \omega_\mu j^\mu$, where

$$j^\mu = \frac{1}{24\pi^2}\epsilon^{\mu\nu\rho\sigma}\,\mathrm{tr}\left[U^{-1}\partial_\nu U U^{-1}\partial_\rho U U^{-1}\partial_\sigma U\right] \quad (6)$$

is the baryonic current. This current is conserved algebraically and can be identified as the topological current connected with the configuration of the unitary field $U(x)$. The total baryonic charge $\int j_0(\mathbf{x},t)d^3x$ then becomes the topological charge, or, the winding number of the unitary field. This model has been extensively studied by three groups in order to see whether the nucleon can be described as a topological soliton.[10,11,12] To predict the low energy properties of the soliton, the following ansatz, known as the hedgehog ansatz, is made : $U(\mathbf{r}) = \exp(i\boldsymbol{\tau}\cdot\hat{\mathbf{r}}\theta(r))$, i.e., $\phi(\mathbf{r}) = \hat{\mathbf{r}}f_\pi\theta(r)$.We refer to $\theta(r)$ as the pion profile function, or, the pion field configuration. What these groups have accomplished is quite significant. They fix all the parameters of the model by studying the meson sector. Then, they proceed to predict the properties of the nucleon as a soliton.[9] All three groups have found that the model describes the low energy properties of the nucleon quite satisfactorily except in one respect. The mass of the soliton comes out consistently too large ($\sim$ 1500 MeV) compared to the actual mass of the nucleon (939 MeV).

Before addressing the problem of large soliton mass, we examine how the topological soliton description of the nucleon gets linked in our work with the high energy elastic scattering.

III. HIGH ENERGY ELASTIC PP AND $\bar{P}P$ SCATTERING AND NUCLEON STRUCTURE

High energy elastic pp and $\bar{p}p$ scattering at the CERN ISR and SPS Collider[13–15] have been analyzed by my collaborators and me over a number of years.[16–18] From our analyses, we have arrived at the following phenomenological description. The nucleon has an inner core and an outer cloud (Fig.1). Elastic scattering at high energy is primarily due to two processes: (1) a glancing collision in which the outer cloud of one nucleon interacts with that of the other giving rise to diffraction scattering; (2) a hard collison in which one nucleon core scatters off the other core via vector meson ω exchange, while their outer clouds overlap and interact independently (Fig. 2). In the small momentum transfer region, diffrac-

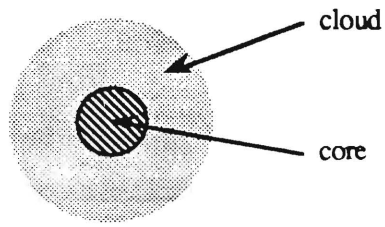

Fig.1. Physical picture of the nucleon from high energy elastic scattering.

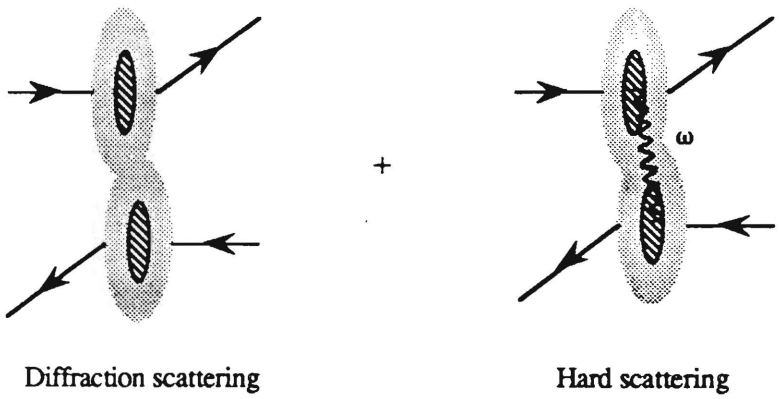

Fig.2. Primary processes responsible for high energy elastic scattering in our description.

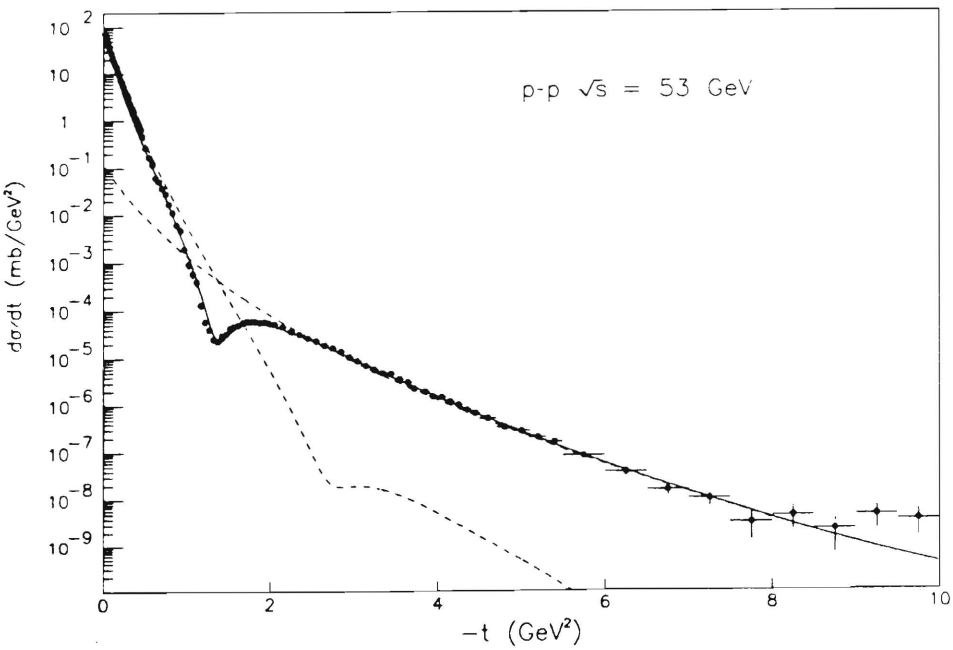

Fig.3. Solid curve represents calculated pp elastic differential cross section at $\sqrt{s} = 53$ GeV [Ref.18]. Differential cross section due to diffraction alone (dot-dashed line) and due to ω-exchange alone (dashed line) are also shown. ISR data from Nagy et al. [Ref.13].

tion dominates. As the momentum transfer increases, the hard scattering takes over. Fig.3 shows our fit to the ISR pp elastic data at $\sqrt{s} = 53$ GeV.[18] The solid curve represents the pp differential cross section calculated by us. The dot-dashed curve represents the differential cross section due to diffraction alone, while the dashed curve represents that due to hard scattering alone, i.e., due to ω exchange. As can be seen, diffraction dominates in the forward direction, but the hard scatterring takes over as the momentum transfer increases.

In our analyses, ω behaves as an elementary spin-1 boson. This can be seen by noticing that the ω exchange amplitude is dominantly of the form:[18]

$$T_H(s,t) \sim e^{i\chi_D(s,b=0)} s \, \frac{F^2(t)}{m_\omega^2 - t} \, .$$

The factor of s originates from the spin 1 of ω. The t-dependence (product of two form factors and the ω propagator) shows that ω exchange corresponds to an elementary meson exchange between two sources. (The factor $e^{i\chi_D(s,b=0)}$ reflects the absorption effect due to diffraction scattering.) From the viewpoint of conventional high energy scattering models, this behavior of ω is hard to understand. In such models, ω is a composite object, and therefore, behaves as a Regge pole. Its contribution away from the forward direction is then negligible. In contrast, the behavior of ω as a spin-1 boson even at high energy is easy to understand in the gauged nonlinear σ-model. This is an effective QCD model in which baryonic charge is topological, and ω is the gauge boson coupled to this charge. As long as the model remains valid, baryonic charge continues to behave as topological charge and ω continues to act as a gauge boson, i.e., as an elementary spin-1 boson. Apparently, this is the behavior we seem to have observed in our phenomenological analyses of elastic scattering. This conclusion, interestingly enough, can be subjected to an important test.[6]

Let us recall the hedgehog ansatz for the soliton model. This ansatz leads to the following expression for the baryonic charge density:

$$\tilde{\rho}(r) = -\frac{1}{2\pi^2} \frac{\sin^2\theta(r)}{r^2} \frac{d\theta(r)}{dr} \, , \tag{7}$$

where $\theta(r)$ is the pion field configuration introduced earlier. If we identify ω as the gauge boson probing the baryonic charge, then from the ωNN form factor found in elastic scattering, we can obtain the baryonic charge density $\tilde{\rho}(r)$, and from $\tilde{\rho}(r)$ we can determine via Eq.(7) the pion field configuration in a nucleon as seen in elastic scattering. This can then be compared with the pion field configuration obtained from the low energy soliton model calculation. The result of such a comparison is shown in Fig.4. The solid curve shows $\theta(r)$ obtained from elastic scattering.The dotted and dashed curves represent the calculations of Meissner et al.[10] in the minimum soliton model (only ω present) and in the complete soliton model (equivalent to having ω, ρ, A_1—all three present). As can be seen, the pion

field configuration from two totally different regimes of physics are not that far apart.

A question that naturally arises in our description of elastic scattering is the following: ω, in our analyses, plays a dominant role, but, ρ and A_1, which are similar low energy vector mesons, play no role. Why does this happen? The nonlinear σ-model again provides an answer. In this model, one introduces ω, ρ and A_1— all on the same footing as gauge bosons. What one finds is that ρ and A_1 couple via isospin vector and axial vector currents, and these currents are dominantly Noether currents.[19] In a path-integral representation, they occur in the usual action. ω, on the other hand, couples via a topological current that occurs in the WZ action. Thus ρ and A_1 couple in a conventional way, and we expect them to behave in a perturbative field theory manner and to Reggeize. Their contributions away from the forward direction would then be negligible. In contrast, ω couples in a nonperturbative way and can continue to act as a gauge boson coupled to the baryonic charge, so that its contribution away from the forward direction can remain significant.

IV. LARGE SOLITON MASS PROBLEM AND GAUGED LINEAR σ-MODEL

We return to the problem of large soliton mass predicted by the gauged nonlinear σ-model. The problem now takes on added importance, since we have found evidence in favor of the model from high energy elastic scattering. To tackle the problem, we go back to the path-integral representation for the gauged linear σ model, extract the WZ action from the fermion measure, and make the approximation that the scalar field $\zeta(x)$ can be replaced by its vacuum value f_π in the meson sector, but not in the quark sector. The partition function then takes the form:

$$\begin{aligned} Tr\left[e^{-iHT}\right] \simeq \frac{1}{\mathcal{N}} \int DU \exp\Big[i\int d^4x \, \frac{1}{4} f_\pi^2 tr[\mathcal{D}_\mu U(\mathcal{D}^\mu U)^\dagger] + i\Gamma_{WZ}[\mathcal{A},\mathcal{B},U]\Big] \\ \times \int D\zeta D\psi^o D\bar{\psi}^o \exp\Big[i\int d^4x \Big(\bar{\psi}_L^o i\gamma^\mu(\partial_\mu + \mathcal{A}_\mu^o)\psi_L^o + \bar{\psi}_R^o i\gamma^\mu(\partial_\mu + \mathcal{B}_\mu^o)\psi_R^o \\ + \frac{1}{2}\partial_\mu\zeta\partial^\mu\zeta - \lambda(\zeta^2 - f_\pi^2)^2 - G\zeta(\bar{\psi}_L^o\psi_R^o + \bar{\psi}_R^o\psi_L^o)\Big)\Big] \; ; \end{aligned} \tag{8}$$

it becomes the product of two functional integrals.The first integral is exactly the same as before (Eq.(5)). It represents the gauged nonlinear σ-model. The second integral shows that, we now have an explicit quark sector where left and right quarks interact via a scalar field, and that this sector can contribute to the total energy of the system. The superscript zero indicates that the fields in the quark sector are all in the unitary gauge (U=1). Further investigation reveals the following.[5]

The nonlinear σ-model, by ignoring interactions in the quark sector, is stating that all the important interactions are in the pseudoscalar meson sector. Quarks are present, but they can be regarded as noninteracting. Therefore, in this model, one has a soliton surrounded by a noninteracting Dirac sea of quarks.On the other hand, if we adopt the linear σ-model and keep the scalar field, then the scalar field mediates an interaction between the left and right quarks, and we have a soliton surrounded by an interacting Dirac sea. What one finds is that if the scalar field has a critical behavior (i.e., zero at small distances, but rises sharply to its vacuum value f_π at some distance $r = R$), then the energy of the interacting Dirac sea can be significantly less than the energy of the noninteracting Dirac sea. The system, consequently, can make a phase transition to the interacting ground state and reduce its total energy or mass by the condensation energy. This mechanism can solve the large soliton mass problem. Furthermore, since it does not affect the pseudoscalar meson sector, all the good predictions of the gauged nonlinear σ-model remain valid. The condensed ground state originating from the interaction between left and right quarks is analogous to a superconducting ground state and can be explicitly written as:

$$|\Psi_o\rangle = \frac{1}{\sqrt{N}} \prod_k \xi_{pk} \xi_{\bar{p}k} \, |\Phi_o\rangle \; ; \tag{9}$$

here, ξ_{pk} and $\xi_{\bar{p}k}$ are the quasi-quark and quasi-antiquark annihilation operators and are related to the usual quark and antiquark annihilation operators by a Bogoliubov transformation. $|\Psi_o\rangle$ is the quasiparticle vacuum, and $|\Phi_o\rangle$ is the particle vacuum. A number of consequences now follow, and let us briefly examine them.

V. CONSEQUENCES OF A CONDENSED GROUND STATE OF CHIRAL QUARKS

1. In our phenomenological description of high energy elastic scattering, the nucleon has an outer cloud. Such a cloud does not occur in the usual low energy soliton model calculations, where the soliton is surrounded by an ordinary vacuum.We now notice that a condensed ground state of chiral quarks (the chiral vacuum) can, in fact, provide the soliton an outer cloud making the low energy picture of the nucleon coincide with its high energy description.

2. Confinement of quarks now becomes automatic.The single-particle excitations of the condensed chiral ground state are quasi-quarks. They exist in the Hilbert space based on the chiral vacuum $|\Psi_o\rangle$, but not in the Hilbert space based on the ordinary vacuum $|\Phi_o\rangle$. Nambu and Jona-Lasinio[20] pointed out a long time ago that the two vacua $|\Psi_o\rangle$ and $|\Phi_o\rangle$ do not overlap, and that there cannot be any transition between the two Hilbert spaces. Therefore, we are not going to see

the quasi-quarks as free particles in the Hilbert space based on $|\Phi_o\rangle$, the ordinary vacuum.

3. If we write the scalar field $\zeta(x) = f_\pi + h(x)$, then the fluctuation field $h(x)$ is the Higgs field in our present context. Its quanta are going to be isoscalar scalar particles, and when two nucleons come close to each other, they will exchange such a Higgs scalar. This Higgs scalar, however, will not be observed as a physical particle or resonance, since it exists only in the condensed chiral vacuum. We can identify it with the σ particle of low energy NN interaction. It has long been known that the attractive intermediate range NN interaction can be attributed to the exchange of an isoscalar scalar particle of mass $\sim$ 500 MeV, which, however, has never been observed experimentally.[21]

4. The condensed ground state $|\Psi_o\rangle$ exhibits spontaneous breakdown of chiral symmetry. If $\hat{Q}_A$ is the conserved axial or chiral charge operator, then it transforms the left and right chiral fields $\hat{\psi}_L$ and $\hat{\psi}_R$ in opposite ways. Consequently, the operator $e^{-i\alpha\hat{Q}_A}$ (α a constant) acting on the vacuum $|\Psi_o\rangle$ transforms it to a new vacuum $|\Psi'_o\rangle$:

$$\begin{aligned} |\Psi'_o\rangle &= e^{-i\alpha\hat{Q}_A}\,|\Psi_o\rangle \\ &= e^{-i\alpha\hat{Q}_A}\frac{1}{\sqrt{N}}\prod_k \xi_{pk}\xi_{\bar{p}k}\,|\Phi_o\rangle \\ &= \frac{1}{\sqrt{N}}\prod_k \xi'_{pk}\xi'_{\bar{p}k}\,|\Phi_o\rangle \quad , \end{aligned} \tag{10}$$

where ξ'_{pk} and $\xi'_{\bar{p}k}$ are the new quasiparticle operators. Since $|\Psi'_o\rangle \neq |\Psi_o\rangle$, we have spontaneous breakdown of chiral symmetry. This is analogous to the spontaneous breakdown of the total charge operator in the BCS superconducting ground state.

VI. CONCLUDING REMARKS

In summary, we see that the gauged linear σ-model brings an explicit quark sector to the gauged nonlinear σ-model. In this sector, left and right quarks interact via a scalar field. A critical behavior of the scalar field can lead to a condensed ground state analogous to a superconducting ground state, and the system can significantly reduce its total energy or mass by the condensation energy. This can solve the large soliton mass problem of the gauged nonlinear σ-model. The nucleon emerges as a topological soliton surrounded by a chiral condensate (Fig.5).[22] The low energy soliton model and the high energy phenomenological description combine to provide this physical picture of the nucleon. If future quantitative calculations bear out the predictions of this combined model, then we would conclude that the elastic data from the CERN ISR and SPS Collider had been pivotal in determining the structure of the nucleon.

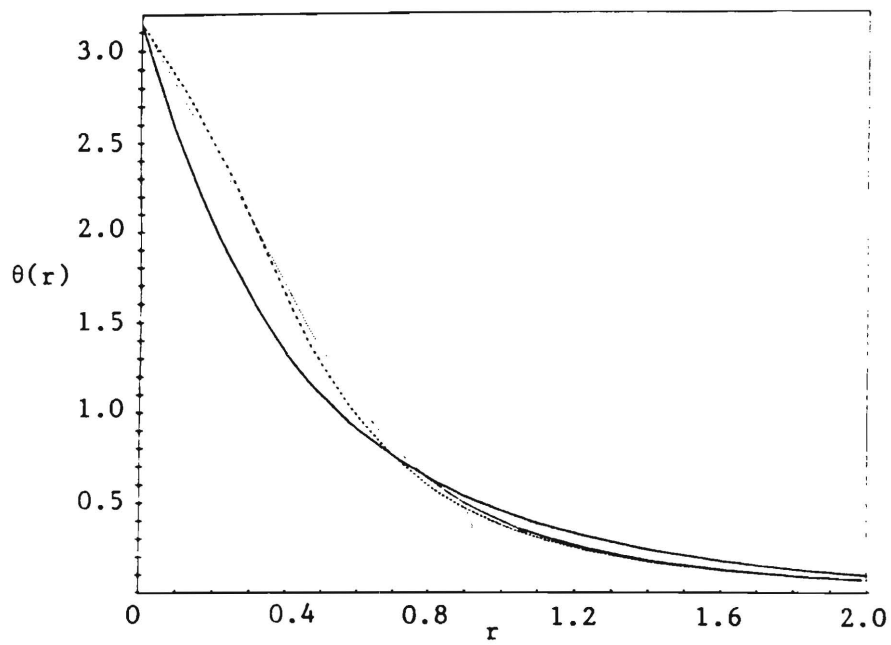

Fig.4. The pion field configuration $\theta(r)$ as a function of r in fermi. The continuous curve represents $\theta(r)$ obtained from high energy elastic scattering. The dotted and the dashed curves represent $\theta(r)$ calculated from low energy in the minimal and in the complete soliton model by Meissner et. al. [Ref.10].

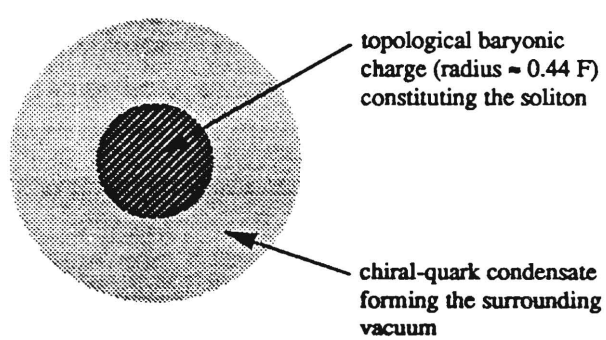

Fig.5 Structure of the nucleon as a topological soliton surrounded by a chiral condensate.

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