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Recovery from Spontaneous Breakdown of Chiral and Gauge Symmetry in Dense Domains of Hadronic Matter

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Abstract

It is speculated that chiral symmetry and unified gauge symmetry of weak and electromagnetic interactions become exact and manifest in regions of high nuclear density.

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DISCLAIMER

A central feature of unified gauge theories of weak and electromagnetic interactions,[1] and of chiral symmetry realization, is the role of the vacuum and its invariance properties. Typically a scalar field multiplet sees a potential whose minimum has a lower symmetry than that of the Lagrangian; the symmetry of the theory is "spontaneously broken" and is realized either with Goldstone-Nambu particles or, with gauge fields present, through the Higgs' mechanism. By choosing the shape of the scalar field potential, a symmetric or asymmetric vacuum may be arranged as the ground state and the physics of one kind of realization (say spontaneously broken, with massive gauge fields) may be smoothly varied into another kind of realization (say exact and manifest, with massless gauge fields) by varying the shape of the potential, that is, by changing the theory.

But is this change of the theory, however appealing, a mere Gedanken experiment? Or can nature itself deform the effective potential, at least in a limited domain of space, and thus physically realize the various modes of symmetry breaking?

Lee and Wick[2] have recently pointed out that a change in an effective scalar field potential may be achieved in a localized region of high nuclear density leading to a stable or metastable "abnormal nuclear state," hereafter denoted ANS. The mechanism, they also note, may alternatively be taken as the basis of a model for binding quarks in a hadron.

Here both possibilities are entertained and it is suggested, in the light of this mechanism, that the chiral symmetry and weak interactions of nucleons in an ANS (or of quarks in a hadron) may be considerably different from those of free nucleons (or free quarks) because a high nuclear (or quark) density induces a vacuum "polarization" which draws the vacuum away from the zero density "direction" to some other domain configuration with profoundly different physical properties.

A consequence is that the ANS may exhibit weak interactions of electromagnetic strength and manifest chiral symmetry exactly with zero effective nucleon and gauge meson masses; and hadron theories which incorporate weak interactions via effectively free constituent quarks, ignoring the binding effects associated with high quark density, may lack central features of the physics.

The possibilities of alternative vacuum realizations of these kinds are as numerous and varied as the many models[1] of unified weak and electromagnetic interactions which have been concocted in the past few years. Here the main idea of perhaps the most radical possibility will be illustrated by studying a hybrid of the original Weinberg lepton theory[3] and the π-σ-N $SU(2)\times SU(2)$ σ-model[4] joined to each other with a cross coupling term[5,6] between the σ-π multiplet and the leptonic Higgs field multiplet. Neglecting nuclear density effects, the model incorporates a conventional picture of chiral symmetry breakdown within the context of a unified weak and electromagnetic gauge theory with a simple extension to hadrons. Details of the model are discussed in Ref.6.

The Lagrangian is

$$\mathcal{L} = \mathcal{L}_W + \mathcal{L}_{WH} + \mathcal{L}_H \tag{1}$$

where

$$\mathcal{L}_W = -\tfrac{1}{4}(\partial_\mu \vec{A}_\nu - \partial_\nu \vec{A}_\mu + g\vec{A}_\mu \times \vec{A}_\nu)^2 - \tfrac{1}{4}(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 - \bar{L}\gamma^\mu D_\mu L$$
$$-\bar{R}\gamma^\mu D_\mu R - \tfrac{1}{2}|D_\mu \varphi|^2 - G_e(\bar{L}\varphi R + \bar{R}\varphi^\dagger L) - M_1{}^2\varphi^\dagger\varphi + h(\varphi^\dagger\varphi)^2, \tag{2}$$

$$\mathcal{L}_{WH} = f(\Sigma_L^\dagger\varphi + \varphi^\dagger\Sigma_L), \tag{3}$$

and

$$\mathcal{L}_H = -\bar{N}_L\gamma^\mu D_\mu N_L - \bar{N}_R\gamma^\mu D_\mu N_R + g_o(\bar{N}_L M^\dagger N_R + \bar{N}_R M N_L) - \tfrac{1}{2}|D_\mu \Sigma_L|^2$$

$$-\lambda_o(\sigma^2 + \vec{\pi}^2 - a^2)^2 \qquad (4)$$

with

$$\varphi = \begin{pmatrix} \varphi^\dagger \\ \varphi^o \end{pmatrix}, \quad \text{left doublet,}$$

$$L = \tfrac{1}{2}(1+\gamma_5)\begin{pmatrix} \nu \\ e \end{pmatrix}, \quad \text{left doublet,}$$

$$R = \tfrac{1}{2}(1-\gamma_5)e, \quad \text{left singlet,}$$

$$\vec{A}_\mu, \ SU(2)_L \text{ gauge triplet,}$$

$$B_\mu, Y_L \text{ gauge singlet,}$$

$$N_L = \tfrac{1}{2}(1+\gamma_5)\begin{pmatrix} p \\ n \end{pmatrix}, \quad \text{left doublet,}$$

$$N_R = \tfrac{1}{2}(1-\gamma_5)\begin{pmatrix} p \\ n \end{pmatrix}, \quad \text{left singlet,}$$

$$M = \sigma + i\vec{\tau}\cdot\pi = \begin{pmatrix} \sigma + i\pi^o & i\sqrt{2}\,\pi^+ \\ i\sqrt{2}\,\pi^- & \sigma - i\pi^o \end{pmatrix},$$

$$\Sigma_L = \begin{pmatrix} -i\sqrt{2}\,\pi^\dagger \\ \sigma + i\pi^o \end{pmatrix}, \quad \text{left doublet,} \qquad (5)$$

where

$$D_\mu = \partial_\mu - ig\vec{t}\cdot\vec{A}_\mu - ig'YB_\mu, \qquad (6)$$

is the $SU(2)_L \times Y$ covariant derivative. $\mathcal{L}_W$ is the Weinberg lepton theory of Ref.3. $\mathcal{L}_H$, apart from terms arising from the covariant derivation is a conventional $SU(2)_R \times SU(2)_L$ σ-model. $\mathcal{L}_{WH}$ is the cross coupling term which preserves the weak $SU(2)_L \times Y$ gauge invariance. In Ref.6 it is shown, in tree approximation, that

$$m_\pi^2 = 4\lambda_o\left(\langle\sigma\rangle^2 - a^2\right) \qquad (7)$$

$$m_\sigma^2 = m_\pi^2 + 8\lambda_o\langle\sigma\rangle^2 \qquad (8)$$

$$M_W = \tfrac{1}{2}g\left(\langle\varphi^o\rangle^2 + \langle\sigma\rangle^2\right)^{\frac{1}{2}}$$

$$M_Z = \tfrac{1}{2}\left(g^2 + g'^2\right)\left(\langle\varphi^o\rangle^2 + \langle\sigma\rangle^2\right)^{\frac{1}{2}} \qquad (9)$$

$$m_N^2 = g_o^2\langle\sigma\rangle^2 \qquad (10)$$

$$\partial^\mu A_\mu^{\ i} = m_\pi^2\langle\sigma\rangle\pi_i + \text{non-pole-terms} \qquad (11)$$

where $\langle\varphi^o\rangle$ is real, $\varphi_1 = \frac{\varphi^o + \varphi^{o\dagger}}{\sqrt{2}}$, and $\langle\sigma\rangle$ and $\langle\varphi_1\rangle = \sqrt{2}\ \langle\varphi^o\rangle$ are solutions of the extremum equations

$$4\lambda_o\sigma(\sigma^2 - a^2) - \sqrt{2}\ \varphi_1 f = 0 \qquad (12)$$

$$M_1^2\varphi_1 - h\ \varphi_1^3 - \sqrt{2}\, f\,\sigma = 0 \qquad (13)$$

arising from minimizing the relevant part of the interaction energy

$$\mathcal{V}(\sigma,\varphi_1) = \lambda_o(\sigma^2 - a^2)^2 - \sqrt{2}\,\sigma\varphi_1 f + \frac{M_1^2\varphi_1^2}{2} - h\left(\frac{\varphi_1^2}{2}\right)^2 \qquad (14)$$

all terms which can not possibly develop vacuum expectation values having been dropped. <u>We shall concentrate on the two interesting cases</u> (1) $h \approx 0$ and (2) $M_1^2 \approx 0$ <u>which lock the chiral breakdown to the gauge symmetry breakdown</u>. Then with the corresponding constraints of Eq.(13) imposed one can define the energy densities

$$U^{(1)}(\sigma) \equiv \mathcal{V}\left(\sigma, \varphi_1 = \frac{\sqrt{2}\,\sigma f}{M_1^2}\right) \qquad (15)$$

$$U^{(2)}(\sigma) \equiv \mathcal{V}\left(\sigma, \varphi_1 = \left(\frac{-\sqrt{2}\, f}{h}\right)^{\frac{1}{3}} \sigma^{\frac{1}{3}}\right) \tag{16}$$

$U^{(1,2)}(\sigma)$ simply project $\mathcal{V}(\sigma,\varphi_1)$ on the surface implied by the constraint of Eq.(13). They are the effective potentials seen by the σ field, with extremum at the solutions of Eq.(12) with (13) imposed. In case (1) this minimum is at

$$\sigma_o^{\,2} = a^2 + \frac{f^2}{M_1^{\,2}\lambda_o}$$

and Eq.(15) can be written

$$U^{(1)}(\sigma) = \frac{m_\sigma^{\,2} - m_\pi^{\,2}}{8}\left(\frac{m_N}{g_o}\right)^2 \left\{\left(\frac{\sigma}{\sigma_o}\right)^2 - 1\right\}^2 \tag{17}$$

where constant terms have been dropped and Eqs.(7), (8), and (10) have been used. Similarly

$$U^{(2)}(\sigma) = \frac{m_\sigma^{\,2} - m_\pi^{\,2}}{8}\left(\frac{m_N}{g_o}\right)^2 \left\{\left[\left(\frac{\sigma}{\sigma_o}\right)^2 - 1 + \delta\right]^2 - 3\delta\left(\frac{\sigma}{\sigma_o}\right)^{\frac{4}{3}}\right\} \tag{18}$$

where

$$\delta = \frac{2m_\pi^{\,2}}{m_\sigma^{\,2} - m_\pi^{\,2}}$$

and in both cases

$$m_N^{\,2} = g_o^{\,2}\sigma_o^{\,2} = g_o^{\,2}\langle\sigma\rangle^2, \tag{19}$$

the effective nuclear mass arising from a non zero vacuum expectation $\sigma_o = \langle\sigma\rangle$, the minimum σ field configuration.

Let us now add to $U^{(1,2)}(\sigma)$ a term describing the effect of high nucleon density, using the incompressible fluid model employed by Lee and Wick[2] and Lee[7]

$$U_N(n,\sigma) = \frac{2}{\pi^2}\int_0^{K_F} k^2 dk\,(k^2 + m_{eff}^2) - n\, m_N \tag{20}$$

where

$$K_F = (3\pi^2 n/2)^{\frac{1}{3}}$$

$$m_{eff}^2 = g_o^{\ 2}\langle\sigma\rangle^2$$

and n is the nucleon density. Then define

$$U_{eff}^{(1,2)}(\sigma,n) \equiv U^{(1,2)}(\sigma) + U_N(n,\sigma),$$

the effective potential seen by the σ field in the presence of nuclear matter of density n. Within such matter the stability position of σ, $\langle\sigma\rangle$, shifts from σ_o, according to the effect of U_N. Figures 1 and 2 are plots of

$$\frac{g_o^{\ 2}}{m_N^{\ 2} m_\sigma^{\ 2}} U^{(1,2)}(\sigma)$$

for various nucleon densities $n = \beta\, n_o$ where n_o is the density of heavy nuclei

$$n_o = \left[\frac{4\pi}{3}\left(1.2\ \text{fm}\right)^3\right]^{-1} = .138\ \text{fm}^{-3}$$

and $(4\pi)^{-1} g_o^{\ 2} = 15$, $m_\sigma = 1.15$ GeV. The similarity of Figs. (1) and (2) indicate that the power with which $\langle\varphi_1\rangle \to 0$ as $\langle\sigma\rangle \to 0$ does not alter the main features of the ANS; indeed, both plots are qualitatively quite similar to the corresponding σ-potential plot of Lee and Wick and Lee, who do not consider the joint σ,φ_1 effect and the implication for chiral and gauge symmetries.

For $\beta = 1$ (normal heavy nucleon density) a metastable ANS is present at $m_{eff}/m_N = 0 = \langle\sigma\rangle/\sigma_o$ which becomes degenerate with the normal state at $\beta \approx 1.1$ and is the stable state for $\beta \geq 1.1$. In such a state, with $\langle\sigma\rangle$ and $\langle\varphi_1\rangle$ both vanishing, chiral symmetry and the weak and electromagnetic gauge symmetry are realized exactly and manifestly, with no spontaneous breakdown. One has effective masses

$$m_\pi^2 = m_\sigma^2 \neq 0$$

$$m_{eff} = 0$$

with the gauge mesons propagating as zero mass particles

$$m_W = m_Z = 0$$

The axial vector current is exactly conserved and SU(2) × SU(2) is realized with mass degenerate multiplets.

Since the gauge symmetry, within these dense domains, is conventionally realized, their weak and electromagnetic interactions will be of equal strength and the usual Wigner-Eckart theorem may be applied to the $SU(2)_L$; for example, the total cross section of neutrinos on lefthanded protons is equal to the total cross section of lefthanded electrons on lefthanded neutrons.

If neutron stars, supernova, or black holes provide the correct density conditions, and ANS are formed with $\langle\sigma\rangle \sim \langle\varphi_1\rangle \sim 0$, then the microscopic astrophysics of such objects will be quite radically different from that of normal matter and structure studies will have to cope with weak interactions as strong as the electromagnetic interaction; such objects may provide a laboratory in which exact symmetry is recovered.

Another possibility is that these states may be found in heavy ion collisions; the signature of such a state, if the $\langle\sigma\rangle \sim \langle\varphi_1\rangle \sim 0$ mechanism is realized, would be a weak decay (say parity violating) at electromagnetic rates.

Finally, the implications for quark theories of hadrons, if the quark density is high enough, include a zero effective quark mass within a hadron with binding dynamics which may involve massless gauge fields and exact realization of unified gauge and chiral symmetry. Work in this direction, similar in spirit to the various field theoretic studies of extended objects, may help provide a dynamical foundation for parton model results.

In summary, what has been suggested here, is that nature displays physically all forms of chiral and unified gauge symmetry realization, from spontaneous breakdown to manifest invariance, this by choosing a branch of extremum solutions in which $\langle\varphi_1\rangle \to 0$ as $\langle\sigma\rangle \to 0$ in regions of high nuclear density; not a necessary choice, but an appealing one for those who like their symmetries exact and manifest, if only in limited domains.

References

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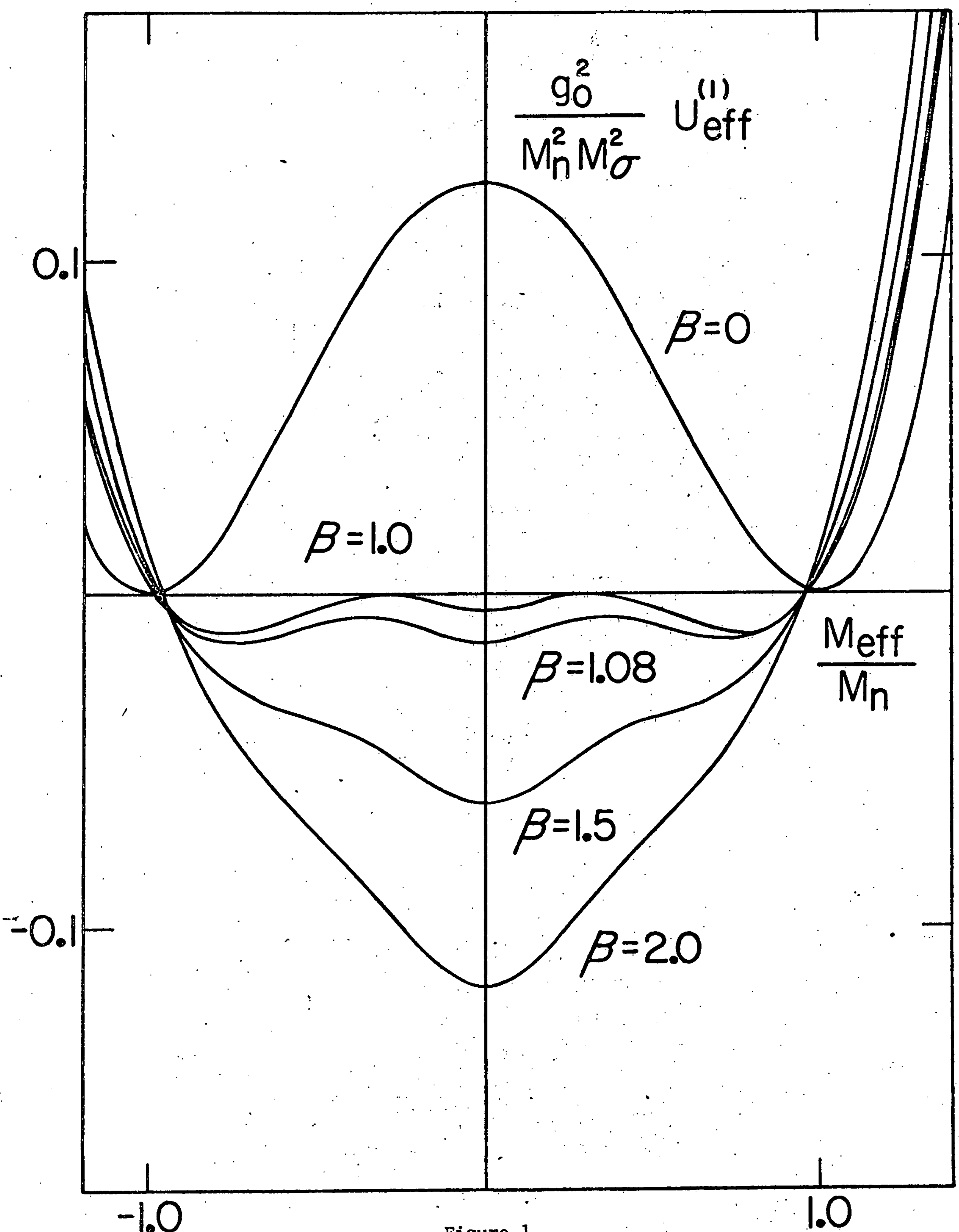
Figure Captions

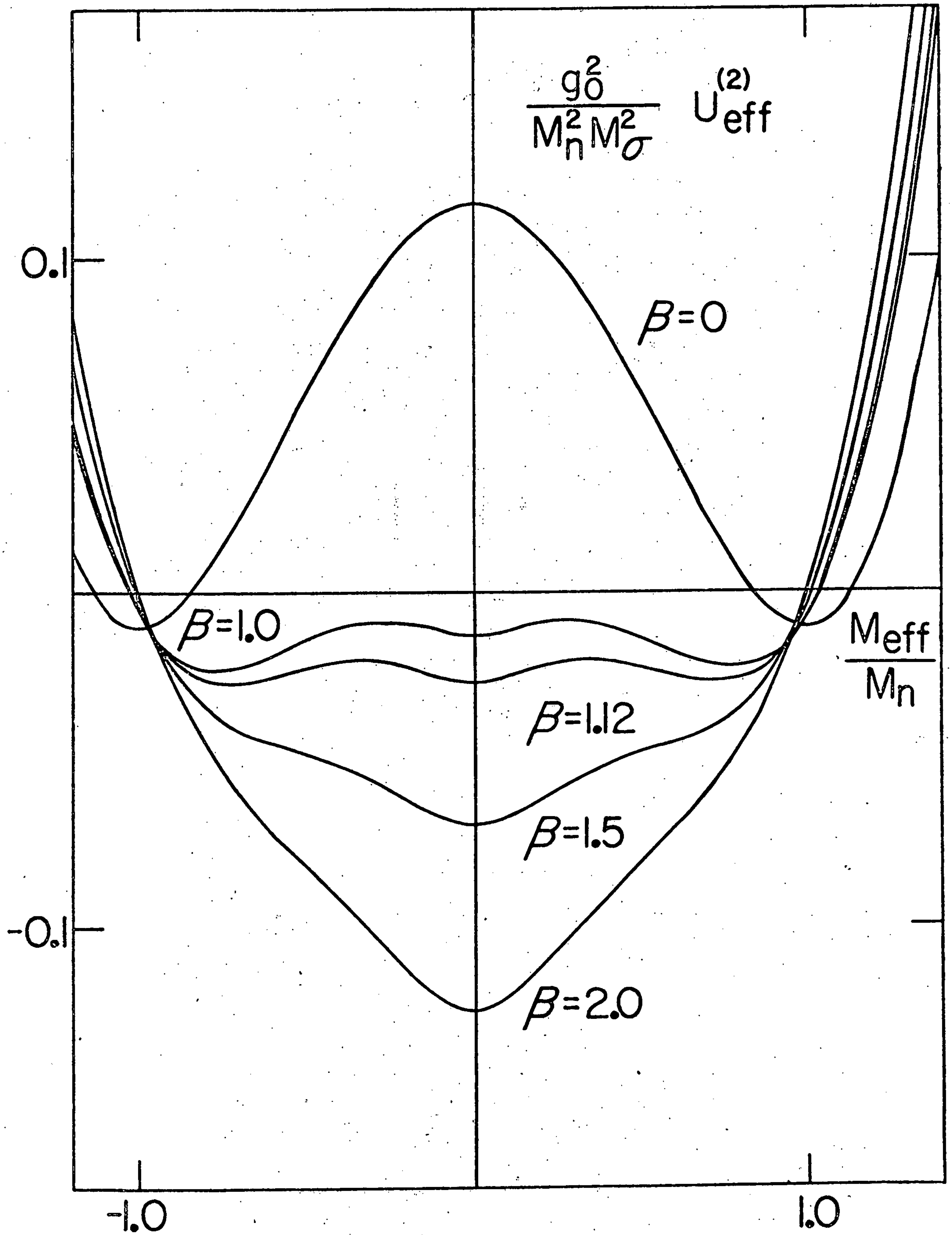
Figure 1. Case (1). The effective potential $U_{eff}^{(1)}(\sigma,n)$ in units of $m_N^2 m_\sigma^2 g_o^{-2}$, seen by the σ field at nuclear density $n=\beta n_o$ as a function of $(\sigma/\sigma_o) = m_{eff}/m_N$ where σ_o is the vacuum expectation at $n=0$. n_o is the nuclear density in heavy nuclei.

Figure 2. Case (2). The effective potential $U_{eff}^{(2)}(\sigma,n)$ in units of $m_N^2 m_\sigma^2 g_o^{-2}$, seen by the σ field at nuclear density $n=\beta n_o$ as a function of $(\sigma/\sigma_o) = m_{eff}/m_N$ where σ_o is the vacuum expectation at $n=0$. n_o is the nuclear density in heavy nuclei.



$\frac{g_0^2}{M_n^2 M_\sigma^2}$ $U_{eff}^{(1)}$
0.1
-0.1
β=0
β=1.0
β=1.08
β=1.5
β=2.0
$\frac{M_{eff}}{M_n}$
-1.0
1.0

Figure 1



$\frac{g_0^2}{M_n^2 M_\sigma^2}$ $U_{eff}^{(2)}$
0.1
-0.1
β=0
β=1.0
β=1.12
β=1.5
β=2.0
$\frac{M_{eff}}{M_n}$
-1.0
1.0

Figure 2