A POSSIBLE NEW FORM OF MATTER
AT HIGH DENSITY

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I.¹ Introduction

In this talk, I would like to discuss some of my recent theoretical speculations, made in collaboration with Gian Carlo Wick. Over the past year, we have tried to investigate the structure of the vacuum. It is through this investigation that the possibilities of vacuum excitation states and abnormal nuclear states have been suggested. Before coming to the main topic, whether or not there may be the possibility of a new form of matter at high density, perhaps I should first digress on questions related to the vacuum.

In physics, one defines the vacuum as the lowest energy state of the system. By definition, it has zero 4-momentum. In most quantum field-theoretic treatments, quite often the vacuum state is used only to enable us to perform the mathematical construct of a Hilbert space. From the vacuum state, we build the one-particle state, then the two-particle state, ...; hopefully, the resulting Hilbert space will eventually resemble our universe. From this approach, different vacuum state means different Hilbert space, and therefore different universe.

Nevertheless, one may ask: What is this vacuum state? Does it have complicated structure? If so, can a part of this structure be changed? Ever since the formulation of relativity, after the downfall of the classical aether concept, one learns that the vacuum is Lorentz invariant. At least, one knows that just running around and changing the reference system won't alter the vacuum. However, Lorentz invariance alone does not insure that the vacuum is necessarily simple. For example, the vacuum can be as complicated as the product or sum of any scalar field or other scalar object at the zero 4-momentum limit:

$$\begin{aligned} \text{vacuum} &\sim \phi^n \text{ or } (\bar{\Psi}\Psi)^m \\ \text{at } k_\mu &= 0 \end{aligned} \tag{1}$$

From Dirac's hole theory, one knows that the vacuum, though Lorentz-invariant, can be rather complicated. That this complicated structure of the vacuum may in part be changeable is suggested by the large variety of broken symmetries, found especially over the past two decades.

If we consider symmetry quantum numbers such as the isospin $\vec{T}$, the strangeness S , the parity P , . . . , we find

$$\frac{d}{dt} \left\{ \begin{array}{c} \vec{T} \\ S \\ P \\ C \\ CP \\ \vdots \end{array} \right\}_{\text{matter}} \neq 0 \quad (2)$$

Aesthetically, this may appear disturbing. Why should nature abandon perfect symmetry? Physically, this also seems mysterious. What happens to these missing quantum numbers? Where do they go to? Can it be that matter alone does not form a closed system? If we also include the vacuum, then perhaps symmetry may be restored:

$$\frac{d}{dt} \left\{ \begin{array}{c} \vec{T} \\ S \\ P \\ C \\ CP \\ \vdots \end{array} \right\}_{\text{matter + vacuum}} = 0 \quad (3)$$

This is, of course, the basic idea underlying the general heading "Spontaneous Symmetry Breaking". On the other hand, unless we have other links connecting matter with the vacuum, how can we be sure that this idea is right, and not merely a tautology?

A way out of this dilemma is to realize that the restriction $k_\mu = 0$ for the vacuum state is only a mathematical idealization. After all, the universe has a finite radius, and k_μ is never strictly zero. So far as the microscopic system of particle physics is concerned, there is little difference between $k_\mu = 0$ and $|k_\mu|$ nearly 0; the latter corresponds to a state that varies only very slowly over a large space-time extension, say of dimension

$$L \gg (\text{microscopic length}) . \quad (4)$$

The typical microscopic length may be ~ 1 fm for a nuclear system, and smaller for a subnuclear system. Thus, it seems worthwhile to search for solutions in a quantum field theory which vary only over a very large space-time volume. As we shall see, excited states of the vacuum may be viewed as examples of such solutions characterized by macroscopic scales.

With these remarks in mind, we now turn to the question of how to generate the vacuum excitation states.

II. Vacuum Excitation

According to (1), the vacuum is $\sim \phi^n$ or $(\bar{\psi}\psi)^m$ at the zero momentum limit. Its excitation can be viewed simply as a change in the expectation value of the corresponding operator ϕ^n or $(\bar{\psi}\psi)^m$ over a large volume. The simplest phenomenological description is to adopt the language of a scalar local field ϕ . The vacuum state is the lowest energy state in such a local field theory. Through the transformation $\phi \rightarrow \phi + \text{constant}$, one may always assume as a convention

$$\langle \text{vac} | \phi | \text{vac} \rangle = 0. \quad (6)$$

A vacuum excitation is represented by an excited state in which the expectation value

$$\begin{aligned} \bar{\phi} \equiv \langle \phi(x) \rangle &\cong \text{constant} \neq 0 && \text{inside } \Omega \\ &= 0 && \text{outside } \Omega \end{aligned} \quad (7)$$

where the linear dimension $\Omega^{\frac{1}{3}}$ is $\gg$ microscopic distance.

How can we produce such changes in $\langle \phi \rangle$? If produced, how can we detect them? The problem is analogous to the formation of "domain structures" in a ferromagnet. We may draw the analog:

In the case of a very large ferromagnet, because its spin interacts linearly with the magnetic field, a domain structure can be created by applying an external magnetic field over a large volume. Furthermore, after domains are created, one may remove the external field; depending on the long-range forces, the surface energy and other factors, such domain structure may persist even after the external magnetic field is removed. Similarly, by applying over a large volume any matter source J which has a linear interaction with $\phi(x)$, one may hope to create a domain structure in $\langle \phi(x) \rangle$. Depending on the dynamical theory, such domains may also remain as physical realities, even after the matter source J is removed. It is important to note that high temperature tends to destroy domain structures. Therefore, throughout this talk, I will concentrate only on zero temperature or low temperature phenomena.

As an illustration, we may consider a local scalar field theory. The Lagrangian density is

$$\mathcal{L}_\phi = -\frac{1}{2} \left(\frac{\partial \phi}{\partial x_\mu} \right)^2 - U(\phi) . \quad (8)$$

If the theory is renormalizable, then U is a fourth order polynomial in ϕ :

$$U(\phi) = \frac{1}{2} a \phi^2 + \frac{1}{3!} b \phi^3 + \frac{1}{4!} c \phi^4 . \quad (9)$$

Since we are interested in the long wavelength limit of the field, the scalar field $\phi(x)$ is used only as a phenomenological description. It can be any kind of 0^+ resonance between particles. The details of its microscopic structure do not concern us, nor is the renormalizability an important factor. All we are interested in is that its zero 4-momentum limit exists; once such a limit exists, then the expectation value $\langle \phi \rangle$ becomes automatically and inextricably connected to the description of the vacuum and the question of vacuum excitations.

In (9), we require the constant $c > 0$ so that U has a lower bound, and $a \geq \frac{1}{2} b^2 / c > 0$ so that the absolute minimum of U is at $\phi = 0$, in accordance with our convention (6). Since we are only interested in a slowly varying field over a large volume, we may neglect both the surface energy and the energy due to $\left(\frac{\partial \phi}{\partial x_\mu}\right)^2$. Therefore, the classical description of $\phi(x)$ should be a reasonably good approximation. The energy density of the field is then simply $U(\phi)$.

1. Constant source

Let us now introduce an external source $J(x)$. As a first example, we consider the simplest case that J is a constant inside a large volume Ω , but zero outside. The energy of the system becomes

$$[U(\phi) + J\phi] \Omega.$$

The following graphs illustrate how in the lowest energy state, inside Ω

$$\bar{\phi} = \langle \phi(x) \rangle \quad (10)$$

can be changed under the influence of a constant J :

In case (ii), as J increases, there is a critical value at which $\bar{\phi}$ makes a sudden jump. Once a domain of an abnormal value of $\bar{\phi}$ is created, if J is subsequently removed, in case (i), $\bar{\phi}$ will return to 0; but in case (ii), depending on the magnitude of the abnormal value produced, even if J is removed $\bar{\phi}$ may not return immediately to 0.

2. Matter source

Let us consider the more realistic case that the source is not a constant, but consists of, say, some physical Fermions, represented by a Dirac field ψ . [Of course, if one wishes, one can also use Bosons instead of Fermions.] For our later discussions on abnormal nuclear states, ψ is the nucleon field; for other purposes such as the "bag" model, ψ can be the quark field. The Lagrangian density becomes

$$\mathcal{L} = \mathcal{L}_\phi - \psi^\dagger \gamma_4 \left[\gamma_\mu \frac{\partial}{\partial x_\mu} + (m + g\phi) \right] \psi \quad (11)$$

where $\mathcal{L}_\phi$ is given by (8), m is the mass of the free Fermion, and g is the coupling constant. An important feature is that if $\bar{\phi}$ is a constant $\neq 0$ over a large volume Ω , then inside Ω the effective mass of the Fermion becomes

$$m_{\text{eff}} = m + g\bar{\phi} \quad (12)$$

Thus, by measuring m_{eff} , one can detect whether there is any change in $\bar{\phi}$. In the above Lagrangian, so far as the Fermion field is concerned, there is an equivalence between

$$m \rightarrow m + \delta m$$

and

$$\phi \rightarrow \phi - \frac{1}{g} \delta m \quad (13)$$

Up to a point, this transformation resembles both the gauge invariance in electrodynamics

$$\begin{aligned}\psi &\rightarrow e^{i\theta} \psi, \\ A_\mu &\rightarrow A_\mu - \frac{1}{e} \frac{\partial \theta}{\partial x_\mu},\end{aligned}\tag{14}$$

and the equivalence between gravitation and acceleration in general relativity. However, in both electromagnetism and gravitation, because of the zero mass of photon and graviton we have exact symmetries. Here, because there is no spin 0 particle of zero mass, transformation (13) is not an exact symmetry in the physical world. [It would, of course, be disastrous if it were, since the inertia mass would then become a non-observable.] As we shall see, this transformation can be used to produce changes in m and ϕ . Thereby, it may lead to a new form of matter at high density.

III. Abnormal Nuclear States

Let the volume Ω be filled with nucleons of density n ,

$$n \equiv \Omega^{-1} \int \langle \psi^\dagger \psi \rangle d^3r . \quad (15)$$

As we shall see, if the nucleon density n is sufficiently high, the system may exist in an "abnormal nuclear state", in which the effective nucleon mass becomes zero, or nearly zero (instead of the free nucleon mass m_N). The simplest way to see why such an abnormal state may develop is to examine the quasi-classical solution.

Let us assume that the nucleons form a Fermi gas of a uniform density and ϕ is a classical field. In such a quasi-classical treatment, the lowest energy state is one in which ϕ is a constant and the Fermi gas is completely degenerate. The corresponding energy density $\mathcal{E}(n)$ is given by

$$\mathcal{E}(n) = U(\phi) + \frac{2}{\pi^2} \int_0^{k_F} (k^2 + m_{\text{eff}}^2)^{\frac{1}{2}} k^2 dk \quad (16)$$

where m_{eff} is the effective mass of the nucleon, related to $m_N \cong 940$ MeV by

$$m_{\text{eff}} = m_N + g\phi , \quad (17)$$

k_F is the Fermi momentum given by

$$k_F = (3\pi^2 n/2)^{\frac{1}{3}} . \quad (18)$$

For simplicity, we assume the nuclear matter to be composed of half protons and half neutrons. We shall also assume Ω to be sufficiently large that the surface energy can be neglected.

Throughout our discussion, we define the normal nuclear state to be one in which $m_{\text{eff}} \cong m_N$ and the abnormal nuclear state to be one in which $m_{\text{eff}} \cong 0$. As n increases, the Fermi-sea contribution to the energy becomes increasingly more important. Thus, independent of the detailed form of $U(\phi)$, in the high density limit one finds

$$\lim_{n \rightarrow \infty} \phi = - (m_N/g) \quad \text{and} \quad \lim_{n \rightarrow \infty} m_{\text{eff}} = 0 ; \quad (19)$$

i.e., the state becomes abnormal. On the other hand, in the low density limit, because of (6), one must have

$$\lim_{n \rightarrow 0} \phi = 0 \quad \text{and} \quad \lim_{n \rightarrow 0} m_{\text{eff}} = m_N ; \quad (20)$$

i.e., the state is normal. As illustrated in the following figure, depending on the parameters in the theory, the transition from the low density "normal" solution ($m_{\text{eff}} \cong m_N$) to the high density "abnormal" solution ($m_{\text{eff}} \cong 0$) may or may not be a continuous one. The mechanism of the transition is similar to that in the case of a constant J discussed before. There is however one important difference: because the nuclear energy depends only on $(m_N + g\phi)^2$; its minimum occurs at $\phi = -m_N/g$. Hence, $\lim_{n \rightarrow \infty} \phi = -m_N/g$ as $n \rightarrow \infty$, while in the case of a constant J one has $\lim_{n \rightarrow \infty} \phi = -\infty$ as $J \rightarrow \infty$.

IV. σ - Model

A particularly interesting example of a discontinuous transition is the well-known σ -model. For a fairly wide range of the parameters, the critical density n_c for the transition is found to be approximately given by

$$n_c \cong 11.6 \left(\frac{m_\sigma}{m_N} \right)^2 \left(\frac{g^2}{4\pi} \right)^{-1} n_0 \quad (21)$$

where

$$n_0^{-1} = \frac{4\pi}{3} (1.2 \text{ fm})^3 ,$$

and m_σ is the σ -meson mass. In the following figure, we map out the region in $\frac{g^2}{4\pi}$ and m_σ for $n_c \leq 2n_0$.

At present, there is no reliable data on either m_σ or g^2 . If we identify σ to be the broad 2π resonance, then

$$m_\sigma \sim 700 \text{ MeV} \quad (22)$$

Extensive calculations have been made in the literature to fit the central attractive part of the nuclear forces from such a 2π exchange. Depending on the assumption of repulsive forces (due to vector meson exchanges), different authors obtained different coupling constants for the $0+$ channel:

$$\frac{g^2}{4\pi} \text{ varies from } 2.5 \text{ to } 15. \quad (23)$$

Even with such a wide range of estimated values, it seems possible that the physical reality may be located at the left of the curve $n_c = 2n_0$. Thus, by doubling the present nucleon density one may produce the abnormal nuclear state.

So far all our discussions are carried out only in the so-called tree approximation. One may wonder about the higher-order loop diagrams. This problem has been systematically investigated recently. All the one-loop and two-loop diagrams have been explicitly calculated; in addition, some of the three- and more-than-three-loop diagrams have also been evaluated by one of my students, Mr. M. Margulies. The next figure gives the result of all one-loop diagrams for $m_\sigma^2 = 1.4 m_N^2$ and $(4\pi)^{-1} g^2 = 5$. For the σ -model

$$U(\sigma) = \frac{m_\sigma^2 m_N^2}{8g^2} \left[1 - \left(\frac{\sigma}{\sigma_0} \right)^2 \right]^2 \quad (24)$$

where $\sigma_0 = m_N/g$ and

$$\bar{\sigma} = \sigma_0 \text{ in the vacuum state.}$$

$\Delta \mathcal{E}_{\text{vac}}$ is the sum of all one-loop diagrams

$$\Delta \mathcal{E}_{\text{vac}} = \text{diagram 1} + \text{diagram 2} + \dots$$

and $\Delta \mathcal{E}'_{\text{vac}}$ includes also

At $\bar{\sigma}^2 < \frac{1}{3} \sigma_0^2$, $\Delta \mathcal{E}_{\text{vac}}$ becomes complex; one must use $\Delta \mathcal{E}'_{\text{vac}}$ instead.

Because the meson-loop diagrams and the nucleon-loop diagrams are always of opposite sign, large cancellations may occur, especially when $m_\sigma^2 \approx 1.4 m_N^2$. This explains why even for $(4\pi)^{-1} g^2 = 5$, $\Delta\mathcal{E}_{\text{vac}}$ and $\Delta\mathcal{E}'_{\text{vac}}$ (calculated for $n = 6n_0$) are both rather small.

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IV. Production and Detection

In order to produce the abnormal nuclear state, we must consider reactions in which (i) a large number of nucleons are involved, so that the surface energy can be neglected, and (ii) the nucleon density can be increased by a significant factor ~ 2 . One is, therefore, led to considerations of high energy collisions between heavy ions, say

where Ab denotes the abnormal nuclear state. If Ab is produced, what are its characteristics? How can it be detected?

Because of the high density involved, the equation of state of the abnormal state depends sensitively on the short-range repulsive force. For definiteness, we may assume it to be a simple hard-sphere repulsion. Let d be the diameter of the hard spheres. In the following figure, we plot the (volume) binding energy per nucleon, $b.e. \equiv \epsilon/n$, versus $d \equiv$ hard-sphere diameter :

$$b.e. \equiv \epsilon/n = \left(\frac{4}{3} \pi n^2 d^3 \right) / n$$

versus

$$d \equiv \text{hard-sphere diameter} :$$

As an example, we may assume $m_\sigma = 700 \text{ MeV}$, $(4\pi)^{-1}g^2 = 5$ and $d = 2/m_\omega \cong .5 \text{ fm}$ where $m_\omega = \omega^0$ meson mass. In this case, per nucleon the kinetic energy is $\sim 490 \text{ MeV}$ and the σ -field energy is $\sim 300 \text{ MeV}$; therefore, one has a binding energy $\sim (940 - 490 - 300) \sim 150 \text{ MeV/nucleon}$ for the abnormal state. The nuclear radius is $\sim (.9 \text{ fm}) A^{\frac{1}{3}}$. If one neglects both the surface energy, and Coulomb energy, the abnormal nuclear state is stable. [For a pure neutron system, the kinetic energy is much higher because of the exclusion principle; in addition, due to ρ^0 -exchange, one expects the repulsive force between neutrons to be greater than that between neutron and proton. Consequently, most likely, the abnormal state of a pure neutron system is unstable.]

A clear signal of reaction (24) is the detection of a stable (or metastable) nucleus of very large baryon number A , say, ~ 400 . For a charge, $Z \sim 200$, the Coulomb energy is not important; therefore, we expect $Z \sim \frac{1}{2} A$. Once an abnormal state is produced, through successive neutron absorption one may hope to increase its A and Z gradually. When Z increases, Coulomb energy becomes important; furthermore, the abnormal state can create e^+e^- pairs. The e^+ will be sent to infinity, but a fair fraction of the e^- will be kept within the abnormal nucleus. As Z increases, the number of e^- also increases. The interplay between the added Fermi energy of e^- and the Coulomb energy may eventually bring the abnormal state to the point of instability when A reaches $\approx 10^4$.

There remains the problem of the production mechanism of the abnormal state. Only a brief discussion will be given, since we are not able to evaluate the various complexities involved. Let us consider reaction (24) at $\sim \frac{1}{2} - 1 \text{ GeV/nucleon}$ in the center-of-mass system.

(i) Penetration problem:

Although the mean free path of a nucleon inside the nuclear matter is relatively short,

because of the high energy, in most of the nucleon-nucleon collisions one expects the two final nucleons to move more or less along their initial general direction, with about 1 or 2 soft mesons created per nucleon-nucleon collision. Thus, provided the impact parameter is small, these two U nuclei should penetrate each other, and thereby increase the nuclear density from its normal value n_0 to $\sim 2n_0$. In addition, there is a large number of soft mesons produced which carry away a large part of the excess energy.

(ii) Response time:

The collision time of (24) is about

$$\tau_{\text{coll}} \sim 10^{-12} \text{ cm}/c$$

where c is the velocity of light (provided that the energy is not very high so that Lorentz contraction can be neglected). The order of magnitude of the response time of the meson field is about

$$\tau_{\text{resp}} \sim \left(\frac{\hbar}{m_\phi c} \right) / c \sim 3 \times 10^{-14} \text{ cm}/c.$$

This follows from the simple Klein-Gordon equation for a scalar field ϕ :

$$-\frac{\partial^2}{\partial x_\mu^2} \phi + m_\phi^2 \phi = -g n_S + \dots$$

where $\dots$ denotes the non-linear part of the meson-meson force and $n_S = \langle \psi^\dagger \gamma_4 \psi \rangle$.

Since

$$\tau_{\text{resp}} \ll \tau_{\text{coll}},$$

as n_S changes, the value of the meson field responds almost adiabatically. When n_S reaches $\sim m_\phi^2 m_N/g$, the meson field becomes $\phi \sim -m_N/g$ which forms the abnormal state.

(iii) Potential barrier:

After the two ^{16}O nuclei penetrate each other, and ϕ develops the abnormal value, there is still the question: would the nucleons keep on moving away from each other, so that the whole system becomes separated? We note that as the meson field assumes the abnormal value $\phi \sim -m_N/g$, the effective nucleon mass becomes zero. The energy of each nucleon is $|\vec{p}|$ inside the nucleus, but $(\vec{p}^2 + m_N^2)^{1/2}$ outside. Thus, there is a potential barrier B holding the nucleons together in the abnormal state. So long as $|\vec{p}|$ is $\lesssim m_N$, most of the nucleons will collide with the barrier and remain inside. Through such collisions, additional mesons will be emitted, and the nucleons may gradually readjust their wave functions to those of the abnormal state.

While an accurate calculation of the production probability is difficult, from the above discussion, one sees that in order to produce Ab , the collision energy of the heavy ions should not be small; otherwise penetration becomes difficult. The collision energy should also be not too high; otherwise the collision time τ_{coll} may become too small because of Lorentz contraction, and the barrier B becomes $\sim \frac{1}{2} |\vec{p}|^{-1} m_p^2$ which may also be too low to hold the nucleons together. The best value may perhaps be ≈ 100 to several hundred MeV/nucleon in the C.M. system.

VI. Remarks

The question whether we live in a "medium" or in a "vacuum" dates back to the beginning of physics. From relativity, we know that the "vacuum" must be Lorentz-invariant. As remarked before, Lorentz invariance by itself does not mean that the "vacuum" is simple. From Dirac's hole theory, one has learned that the vacuum, though Lorentz-invariant, can be rather complicated. However, so long as all of its properties cannot be changed, so long as, e.g., the value of vacuum polarization cannot be modified, then it is purely a question of semantics whether the vacuum should be called a medium or not.

What we try to suggest is that if we do indeed live in a medium, then there should be ways through which we may change the properties of that medium.

It may be worthwhile to emphasize once again how limited have been our experiences in either nuclear physics or particle physics. So far almost all our nuclear physics experiments have been restricted to nuclei at a constant density. We have never really ventured out to study nuclear physics at any densities other than the normal one. Likewise, in particle physics, we have a similar tradition in specializing along a fairly narrow direction. Take, for example, high energy physics. Hitherto, we have concentrated only on experiments in which we distribute a higher and higher amount of energy into a region with smaller and smaller dimensions. In order to study the question of "vacuum", or the possibility of the abnormal states, we must turn to a different direction; we should investigate some "bulk" phenomena by distributing high energy or high nucleon density over a relatively large volume. The fact that such directions have never been explored should, by itself, serve as an incentive for doing such experiments. As we have discussed, there are possibilities that abnormal states may be created, in which

the nucleon mass may be very different from its normal value. It is conceivable that inside the volume of the abnormal state, some of the symmetry properties may become changed, or even that the usual roles of strong and weak interactions may become altered. If indeed the properties of the "vacuum" can be transformed, we may eventually be led to some even more striking consequences than those that have been discussed in this lecture.

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