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HOT METASTABLE STATE OF ABNORMAL MATTER
IN RELATIVISTIC NUCLEAR FIELD THEORY

N.K. Glendenning

December 1986

NUCLEAR SCIENCE DIVISION

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Abstract

Because of their non-linearity, the field equations of relativistic nuclear field theory admit of additional solutions besides the normal state of matter. One of these is a finite-temperature abnormal phase. Over a narrow range in temperature, matter can exist in the abnormal phase at zero pressure. This is a hot metastable state, for which there is a barrier against decay, because the field configuration is different than in the normal state, the baryon masses are far removed from their vacuum masses, there is an abundance of pairs also far removed from their vacuum masses, and a correspondingly high entropy. The abundance of baryon-antibaryon pairs is the glue that holds this matter together. The signals associated with this novel state are quite unusual. A fragment of such matter will cool by emitting a spectrum of black-body radiation, consisting principally of photons, lepton pairs and pions, rather than by baryon emission, because the latter are far removed from their vacuum masses. If produced at the upper end of its temperature range, a large fraction of the original energy, more than half in the examples studied here, is radiated in this way. The baryons and light elements produced in the eventual decay, after the abnormal matter has cooled to a domain where its pressure becomes positive, will account for only a fraction of the original energy. The energy domain of this state depends sensitively on the coupling constants, and within a reasonable range as determined by nuclear matter properties, can lie in the range of GeV to tens of GeV per nucleon.

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1 Introduction

Relativistic nuclear field theory provides a good description of the bulk properties of nuclear matter as well as a large number of single-particle properties of finite nuclei [1]. With appropriate extensions the theory can be used to study matter away from the normal state, matter that is under extreme conditions of temperature or density, such as is expected to be produced in relativistic nuclear collisions, and as formed in the collapse of a star just prior to the bounce that produces the supernova, and as exists in the cores of the neutron stars into which the remaining matter of the star subsides.

In addition to the solution corresponding to the normal state of matter, the theory has the possibility of possessing additional solutions that correspond to different field configurations [2]. The reason for this is that the field equation for the scalar field is non-linear, which is intrinsic to the Yukawa coupling of Fermions to a scalar field. In earlier work [3,4,5,6,7], a second solution was found at high temperature corresponding to a phase of matter that is characterized by low baryon effective masses and a high density of baryon-antibaryon pairs. Here we point out that under certain conditions there exists a temperature range over which matter in this abnormal phase can have zero pressure. In that case it is mechanically stable. Because of its internal properties, including its field configuration, there is a barrier against its decay to the normal state. An object made of such matter is therefore metastable. However since it is hot it will radiate. Because the baryons are far removed from their vacuum masses, there is a barrier against their emission. Instead, a fragment of such matter will cool by black-body radiation until it reaches a domain were its pressure becomes positive. It will then disassemble into baryons and perhaps lighter nuclei. It will evolve in density during its radiation era from near nuclear to super-nuclear densities.

The energy range in which this state exists depends sensitively on the coupling constants, and within their reasonable range as determined by nuclear matter properties, can lie in the range of GeV to tens of GeV per nucleon.

In the following sections we will outline the derivation of the equations necessary to an understanding of the abnormal phase. Then we contrast the gas-liquid and the normal-abnormal phase transitions, and discuss the characteristics of the abnormal phase. Its physical origin is described. The existence of a solution corresponding to a metastable state is observed and the decay properties and signals of a fragment of such matter are described. These are contrasted with the decay of compressed normal matter and a quark-gluon plasma.

2 Theory

2.1 Lagrangian and Field Equations

As we shall discuss latter, abnormal solutions such as outlined above may exist for any theory having a scalar field coupled to baryons. However, the solution has special interest for the theory that can describe known nuclear properties. This is the scalar-vector-isovector theory [8,9,10,11], possibly augmented with ϕ^4 interactions [12]. For hot dense matter, the theory should be extended in several ways. First it is evident that in such matter, nucleons at the top of the Fermi sea have greater energy than the masses of other baryon species, including hyperons. So the excited nucleons, deltas and hyperons should also be incorporated. In hot matter, thermal mesons will also be produced. We made these generalizations in earlier studies. In non-strange matter the hyperons must appear as particle-antiparticle pairs. This would be a constraint on the solution in the case that the matter has a very short lifetime. A metastable object can develop a net strangeness with the decay of the associated kaons if its lifetime exceeds the weak interaction time scale. This will lead to a still more stable configuration.

The extended Lagrangian is [3,4,5]

$$\begin{aligned}\mathcal{L} &= \sum_B \overline{\psi}_B\left[i\gamma_\mu(\partial^\mu + ig_{\omega B}\omega^\mu) - (m_B - g_{\sigma B}\sigma)\right]\psi_B \\ &\qquad - \tfrac{1}{3}bm_n(g_\sigma\sigma)^3 - \tfrac{1}{4}c(g_\sigma\sigma)^4 + \mathcal{L}^0_\sigma + \mathcal{L}^0_\omega + \cdots \qquad (1)\end{aligned}$$

Here ψ_B denotes a baryon spinor and the sum is over all of the baryon families, N, $\Delta, \Lambda, \Sigma, \Xi$... and their excited states until the solution converges over the range of densities considered. The $\mathcal{L}^0$ are the free Lagrangians of the mesons. The σ- and ω-mesons are Yukawa coupled to the baryon scalar density and vector current, and the ρ-meson is coupled to the total isospin current which however vanishes in symmetric matter. The other mesons are included as free thermal bosons and their Lagrangians are represented by the ellipsis.

The field equations for static uniform matter in the mean field approximation are, for the mesons,

$$m_\sigma^2 \sigma = -bm_n(g_\sigma\sigma)^2 - c(g_\sigma\sigma)^3 + \sum_B g_{\sigma B} < \bar{\psi}\psi > \quad (2)$$

$$m_\omega^2 \omega_0 = \sum_B g_{\omega B} < \psi^\dagger\psi > \quad (3)$$

$$m_\omega^2 \omega_k = \sum_B g_{\omega B} < \bar{\psi}\gamma_k\psi > = 0 \implies \omega_k \equiv 0 \quad (4)$$

and for the baryons,

$$\left[\gamma^\mu(p_\mu - g_{\omega B}\omega_\mu) - (m_B - g_{\sigma B}\sigma)\right]\psi_B = 0. \quad (5)$$

The baryon eigenvalues of momentum, p, for particle and antiparticle are,

$$\epsilon_B(p) = E_B(p) + g_{\omega B}\omega_0 \quad (6)$$

$$\bar{\epsilon}_B(p) = E_B(p) - g_{\omega B}\omega_0, \quad (7)$$

with

$$E_B(p) = \sqrt{p^2 + (m_B - g_{\sigma B}\sigma)^2}. \quad (8)$$

The source of attraction in the theory is the scalar field which reduces the masses of the baryons in the medium. We denote the effective mass of species B by,

$$M_B = m_B - g_{\sigma B}\sigma. \quad (9)$$

The value of the nucleon effective mass at saturation density will be denoted by m^* and the vacuum mass by m. In baryon matter, as contrasted with antibaryon matter, the vector field, eq. (3), is positive, being driven by the net baryon density. Therefore the vector field shifts the baryon eigenvalues to higher energy, in baryon rich matter, as seen in eq. (6), and consequently is a source of repulsion. The effect on the antibaryons is opposite, eq. (7). When this shift is summed over the baryon spectrum, the result is a repulsive energy that is quadratic in the density.

The energy density and pressure can be found from the stress-energy tensor whose mean value is,

$$\bar{\mathcal{T}}_{\mu\nu} = -g_{\mu\nu}\bar{\mathcal{L}} + < \bar{\psi}\gamma_\mu p_\nu \psi > \tag{10}$$

Thus,

$$\epsilon \equiv \mathcal{H} = -\bar{\mathcal{L}} + < \bar{\psi}\gamma_0 p_0 \psi > \tag{11}$$

$$p = \bar{\mathcal{L}} + \frac{1}{3} < \bar{\psi}\gamma \cdot \mathbf{p}\psi > \tag{12}$$

where $\bar{\mathcal{L}}$ is the Lagrangian evaluated for the field values that satisfy the field equations,

$$\bar{\mathcal{L}} = \tfrac{1}{2}m_\omega^2\omega_0^2 - \tfrac{1}{2}m_\sigma^2\sigma^2 - \tfrac{1}{3}bm_n(g_\sigma\sigma)^3 - \tfrac{1}{4}c(g_\sigma\sigma)^4, \tag{13}$$

and $g_{\mu\nu}$ is the diagonal metric tensor

$$g_{\mu\nu} = (1, -1, -1, -1). \tag{14}$$

2.2 Finite Temperature

From statistical mechanics, the grand partition function,

$$Z = Tr[e^{-(\hat{H}-\mu\hat{N})/T}], \tag{15}$$

is related to the thermodynamic potential for volume V by,

$$\Omega = pV = -T \ln Z, \tag{16}$$

where $\hat{H}$ is the Hamiltonian operator, $\hat{N}$ is the number operator and μ, the chemical potential. In the present case of a many component system of baryons of species B $(= N, \Delta, \Lambda, \ldots)$, their excited states and antibaryons, $\bar{B}$ and particle momenta $\mathbf{p}$, the following replacement should be made,

$$\mu\hat{N} \longrightarrow \sum_{B,\mathbf{p}}(\mu_B\hat{N}_{B\mathbf{p}} + \bar{\mu}_B\hat{\bar{N}}_{B\mathbf{p}}). \tag{17}$$

Here μ_B and $\bar{\mu}_B$ denote the chemical potential for baryon and antibaryon of species B.

For the theory defined above, we can write the Hamiltonian operator as,

$$\begin{aligned}\hat{H} &= \int_V d^3x\, \hat{\mathcal{H}} \\ &= -\mathcal{L}V + \sum_{B,\mathbf{p}} \left\{ \epsilon_B(p)\hat{N}_{B\mathbf{p}} + \bar{\epsilon}_B(p)\hat{\bar{N}}_{B\mathbf{p}} \right\} + \sum_{M,\mathbf{p}} \epsilon_M(p)\hat{N}_{M\mathbf{p}} \,. \end{aligned} \tag{18}$$

We take into account mesons additional to those that interact with the baryons in eq. (1), by including them as thermal bosons. They are represented in the Hamiltonian by the last term in eq. (18), where M is summed over them and $\epsilon_M(p)$ denotes their energy.

$$\epsilon_M(p) = \sqrt{p^2 + m_M^2}. \tag{19}$$

The trace in eq. (14) can now be evaluated, taking into account the Fermi and Bose statistics of baryons and mesons.

$$\begin{aligned}\ln\, Z &= \mathcal{L}V/T + \sum_{B,\mathbf{p}} \left\{ \ln[1 + e^{-(\epsilon_B(p)-\mu_B)/T}] + \ln[1 + e^{-(\bar{\epsilon}_B(p)-\bar{\mu}_B)/T}] \right\} \\ &\quad - \sum_{M,\mathbf{p}} \ln[1 - e^{-\epsilon_M(p)/T}]. \end{aligned} \tag{20}$$

The expectation values of the number operators in eq. (18) are found now from,

$$< N >= T\Big(\frac{\partial \ln Z}{\partial \mu}\Big)_{T,V} \tag{21}$$

namely,

$$n_B(p) = \frac{1}{1 + \exp[(\epsilon_B(p) - \mu_B)/T]}\,, \tag{22}$$

with a similar expression for antibaryons. The meson distribution is of course,

$$n_M(p) = \frac{1}{\exp[\epsilon_M(p)/T] - 1}. \tag{23}$$

The sums in eq. (19) can be converted to integrals in the usual way,

$$\sum_{\mathbf{p}} \longrightarrow \frac{(2J_B + 1)(2I_B + 1)}{2\pi^2} \int_0^\infty p^2 dp. \tag{24}$$

Thus we obtain,

$$\begin{aligned}\ln Z &= \mathcal{L}V/T + \sum_B \frac{(2J_B+1)(2I_B+1)}{2\pi^2} \\ &\int_0^\infty p^2 dp \Big\{ \ln[1+e^{-(\epsilon_B(p)-\mu_B)/T}] + \ln[1+e^{-(\bar{\epsilon}_B(p)-\bar{\mu}_B)/T}] \Big\} \\ &- \sum_M \frac{(2J_M+1)(2I_M+1)}{2\pi^2} \int_0^\infty p^2 dp \, \ln[1-e^{-\epsilon_M(p)/T}]. \end{aligned} \tag{25}$$

The finite temperature field equations for σ and ω_0 can now be found as the condition that the thermodynamic potential is minimized at constant V, T and μ,

$$\frac{\delta\Omega}{\delta\sigma} = \frac{\delta\Omega}{\delta\omega_0} = 0, \tag{26}$$

namely

$$\begin{aligned} m_\sigma^2 \sigma &= -\, b m_n (g_\sigma \sigma)^2 - c(g_\sigma\sigma)^3 \\ &+ \sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} g_{\sigma B} \\ &\int_0^\infty \frac{m_B - g_{\sigma B}\sigma}{\sqrt{p^2 + (m_B - g_{\sigma B}\sigma)^2}} [n_B(p) + \bar{n}_B(p)]\, p^2 dp \end{aligned} \tag{27}$$

$$\begin{aligned} m_\omega^2 \omega_0 &= \sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} g_{\omega B} \\ &\int_0^\infty [n_B(p) - \bar{n}_B(p)]\, p^2 dp. \end{aligned} \tag{28}$$

Notice that σ occurs on the right side of both equations, since it is contained in the eigenvalue appearing in the Fermi-Dirac functions. This is a highly non-linear equation, since if expanded in a power series in σ, it is of infinite order. The non-linearity is intrinsic to the Yukawa coupling of Fermions to a scalar field. The potential terms with coefficients b and c also introduce a non-linearity but they are not essential to the phenomena that are discussed in this paper.

Notice also that the scalar field is driven by the total (scalar) density of baryons and antibaryons, while the vector field is driven only by the net baryon density,

$$\rho = \sum_B \frac{(2J_B+1)(2I_B+1)}{2\pi^2} \int_0^\infty [n_B(p) - \bar{n}_B(p)]\, p^2 dp, \tag{29}$$

The baryon chemical potential for species B (= N,$\Delta, \Lambda \ldots$), is denoted by μ_B. For antibaryons, $\bar{\mu}_B = -\mu_B$. For non-strange baryons and hyperons, the chemical potentials obey,

$$\mu_B = \mu, \quad \text{for non-strange baryons} \tag{30}$$

$$\mu_B = g_{\omega B}\omega_0, \quad \text{for strange baryons.} \tag{31}$$

The latter condition assures that the hyperons occur in particle-antiparticle pairs, as can be seen by reference to eq. (6,7,22).

When the solutions to the field equations have been found, the pressure and energy density can be calculated. An expression for the pressure can be obtained from eq. (16) and an integration by parts of eq. (25). The energy density can be found from eq. (18). The results are,

$$\begin{aligned} p = & -\tfrac{1}{3}bm_n(g_\sigma\sigma)^3 - \tfrac{1}{4}c(g_\sigma\sigma)^4 - \tfrac{1}{2}m_\sigma^2\sigma^2 + \tfrac{1}{2}m_\omega^2\omega_0^2 \\ & + \frac{1}{3}\sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} \\ & \quad \int_0^\infty \frac{p^2}{\sqrt{p^2+(m_B-g_{\sigma B}\sigma)^2}}\left[n_B(p)+\bar{n}_B(p)\right]p^2dp \\ & + \frac{1}{3}\sum_M \frac{(2I_M+1)(2J_M+1)}{2\pi^2}\int_0^\infty \frac{p^2}{\sqrt{p^2+m_M^2}} n_M(p)p^2dp \end{aligned} \tag{32}$$

$$\begin{aligned} \epsilon = & \tfrac{1}{3}bm_n(g_\sigma\sigma)^3 + \tfrac{1}{4}c(g_\sigma\sigma)^4 + \tfrac{1}{2}m_\sigma^2\sigma^2 + \tfrac{1}{2}m_\omega^2\omega_0^2 \\ & + \sum_B \frac{(2I_B+1)(2J_B+1)}{2\pi^2} \\ & \quad \int_0^\infty \sqrt{p^2+(m_B-g_{\sigma B}\sigma)^2}\left[n_B(p)+\bar{n}_B(p)\right]p^2dp \\ & + \sum_M \frac{(2I_M+1)(2J_M+1)}{2\pi^2}\int_0^\infty \sqrt{p^2+m_M^2}\, n_M(p)p^2dp. \end{aligned} \tag{33}$$

In these equations, σ and ω_0 denote the mean values of the scalar meson, and the time-like component of the ω-meson. The space-like components vanish in isotropic matter. The sum on B is over the baryon species, and the one on M is over the thermal mesons, of which the pion triplet is most important.

We notice that if the sign of the chemical potential, μ, is everywhere changed, the sign of the baryon density and ω_0 is changed. However the scalar field equation remains unchanged as do p and ϵ. Therefore the solutions are reflection symmetric between positive and negative net baryon density (matter and anti-matter). This may be kept in mind when viewing graphs of properties as a function of density.

2.3 Coupling Constants

The four coupling constants in the theory, g_σ/m_σ, g_ω/m_ω, b, c, are chosen so that the theory possesses the bulk properties of uniform symmetric matter, B/A = 15.95 MeV, saturation density $\rho = 0.145 fm^{-3}$ [13], the compression modulus $K = 240$ MeV, and the nucleon effective mass *at saturation*. This quantity is not well known, so we shall use two values, $m^\star/m = 0.8$ and 0.75. The value of the compression modulus is consistent with the analysis of the giant monopole resonance [14], with the droplet model of atomic masses [13], and with the charge-distribution differences of heavy isotopes [15]. It is also consistent with known neutron star masses [16]. Of course nuclear matter properties do not determine the hyperon couplings. For simplicity we have assumed universal coupling. All well established baryons up to mass 1775 MeV were included in the calculation, and the mesons that were included in the thermal ensemble were similarly summed to convergence, although for the temperatures covered in this study, the pion is the most important.

3 High Temperature Properties

3.1 Gas-Liquid and Normal-Abnormal Phase Transitions

At low temperature and density, the theory possess the gas-liquid phase transition [17,18], which for purposes of comparison with the high-temperature phase transition, is illustrated in Fig. 1 and 2. The crossing point on the pressure-chemical potential isotherms exhibits graphically the Gibbs condition of phase equilibrium, namely, that the pressure, temperature and chemical potentials are equal in the two phases. In the region in temperature and density above the gas-liquid phase transition, one might have expected that pressure and energy density would be monotonic increasing functions of baryon density. However, as shown on a high-temperature isotherm in Fig. 3, this is not so. The pressure reaches a maximum on the normal branch, and there is a second state of equal pressure at higher density. Viewed as

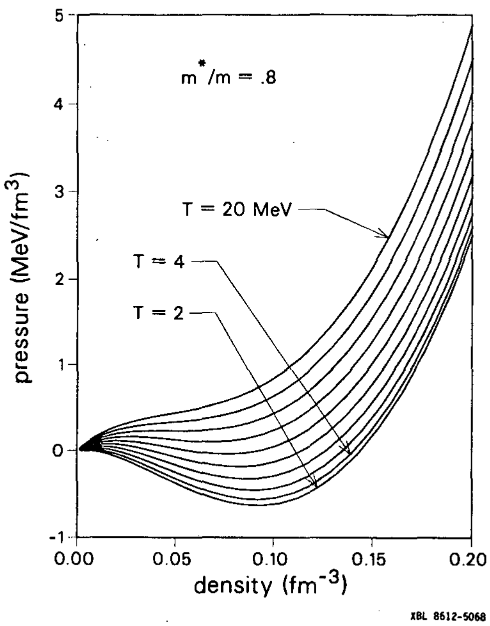

Figure 1: Isotherms of pressure as a function of net baryon density in the region of the gas-liquid phase transition.

a function of density, both pressure and density are three valued functions of each other, in contrast to the gas-liquid phase region, where only the pressure is three-valued. However, as shown schematically in Fig. 4, when pressure is viewed as a function of the scalar density, it is again only the pressure that is three valued. The scalar field amplitude rather than density is the order parameter for the high-temperature phase transition.

The existence of a region in which there are three solutions at the same density was noticed some years ago [3] but was not fully explored. The critical region is shown on energy isotherms in Fig. 5 where for $165 \geq T \geq 146$ MeV three solutions exist at the same density. Two of the three solutions are, in the case of uniform matter, local minima in the thermodynamic potential regarded as a function of σ and the other is a maximum and therefore an unstable solution (the middle one). We have not examined stability with respect to space-varying fields. However, we cannot find any source of attraction in the theory arising from a spatial variation that would more than compensate for the corresponding kinetic energy, so we believe that the minima found for uniform matter are true minima of the thermodynamic potential. We shall refer to the lower energy branch as the normal solution and the higher energy branch as the abnormal one. Its properties will be discussed below.

For temperatures below $T_n \approx 146$ MeV only the normal phase exists,

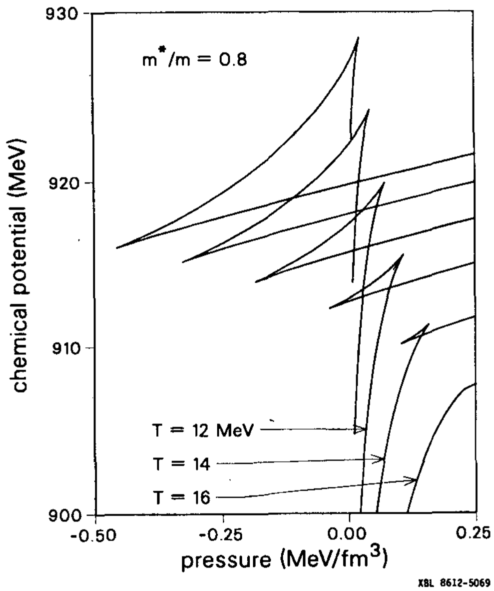

Figure 2: Pressure as a function of chemical potential in the vicinity of the critical point of the gas-liquid phase. The crossing point on each isotherm is the condition of phase equilibrium.

and it spans the entire range of density. Above this critical temperature, a first order phase transition sets in, and the normal phase is confined to a finite range of density beginning at zero. The range of the normal state shrinks with increasing temperature, so that above $T_a \approx 165$ MeV only the abnormal phase exists, and it spans the entire range of density. Between the above two critical temperatures matter can exist in either the normal or abnormal state.

We shall see later that the domain in temperature and density of the critical region of the normal-abnormal phase transition may move down dramatically, depending on the choice made for the nucleon effective mass at saturation, which for the above discussion is $m^*/m = 0.8$.

3.2 Origin of the Abnormal State

The abnormal phase discussed here is a finite temperature one, unlike the density isomer discussed by Lee and Wick [19]. They have in common low baryon effective masses. However the variable that drives the effective mass is quite different. Here it is temperature that drives the phase transition in the following way. At finite temperatures baryon-antibaryon pairs can be created because they have the quantum numbers of the vacuum. Baryon and antibaryon act differently on the scalar and vector fields. Being scalar, the scalar field is driven equally by baryon and antibaryon, as seen in eq. (27).

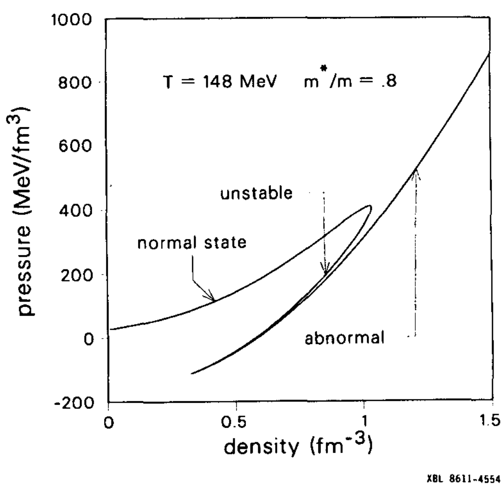

Figure 3: Pressure as a function of density on the isotherm T = 148 MeV., showing details of the normal, abnormal and unstable solution. Note that on the unstable solution the pressure is zero at a density of about $0.6\,fm^{-3}$.

The pairs therefore increase the scalar field strength, reducing their effective mass, thereby making it energetically possible to accommodate even more pairs. Thus an instability can arise, leading to a first order phase transition. The vector field, on the other hand, being vector, *can* distinguish baryon and antibaryon, and *pairs* have null effect on the vector field, eq. (28).

The above reasons for the existence of an abnormal phase in theories in which Fermions are coupled to a scalar field are not particular to the precise form of the Lagrangian, since the non-linearity in the scalar field equation that admits the possibility of multiple solutions is intrinsic to such theories [5].

In contrast to the above discussion, for the Lee-Wick density isomer it is the density of baryons that drives the effective mass, but that also drives the vector repulsion. Such an isomer has not been found in the present scalar-vector theory for that reason. Nevertheless the theory possesses the finite temperature abnormal phase. One can see generally, by the nature of the driving mechanism, that the finite temperature phase transition to abnormal matter can exist in theories for which the Lee-Wick density isomer does not.

The principal characteristics of the abnormal phase, are the large scalar field, leading to low effective baryon masses, the high abundance of baryon-antibaryon pairs, and the corresponding high entropy per net baryon. As can be seen in Fig. 6, the scalar field strength has low to moderate values on the normal branch, but high values, near or above the nucleon mass on

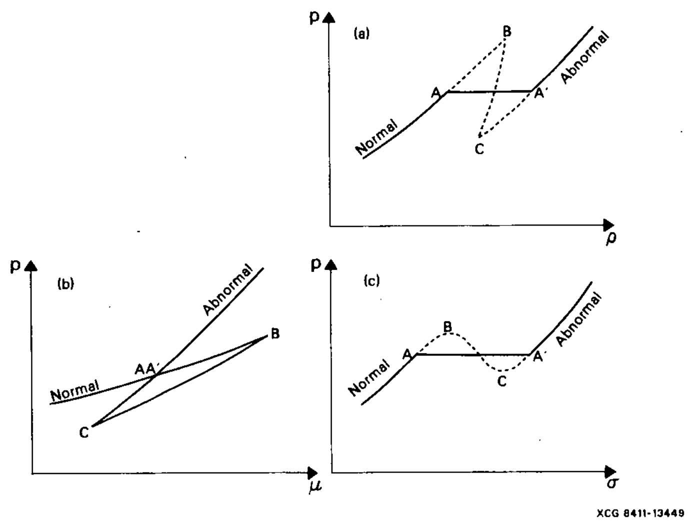

Figure 4: Schematic showing pressure as a function of density (a) and of scalar field strength (c). Part (b) shows the phase equilibrium point, AA'.

the abnormal branch. The entropy per baryon is shown on an isotherm in Fig. 7. Here we see that the entropy is much larger in the abnormal than the normal phase at the same density, or that the normal phase has the same entropy as the abnormal only when the former is highly dispersed.

3.3 Metastable Bodies

In a narrow range of temperatures, 147 MeV $< T <$ 152 MeV, the pressure becomes zero over a segment of the abnormal branch. This is shown for $T =$ 148 MeV in Fig. 3. The point of zero pressure on the abnormal branch is mechanically stable. Clearly the source of binding energy is the high abundance of baryon-antibaryon pairs that are characteristic of the abnormal phase. They reduce the effective masses while having null effect on the repulsive vector field.

The locus of these zero pressure points on the abnormal branch is finite in extension and is shown in the energy-temperature plane in Fig. 8. Let us suppose that in a high energy collision between heavy nuclei a piece of matter has been produced in this state. The energy required to produce the abnormal state at $T =$ 151 MeV in the zero pressure configuration is 22 GeV per nucleon in the center of mass system. Having zero pressure, it is mechanically stable. The surface tension and Coulomb force may actually

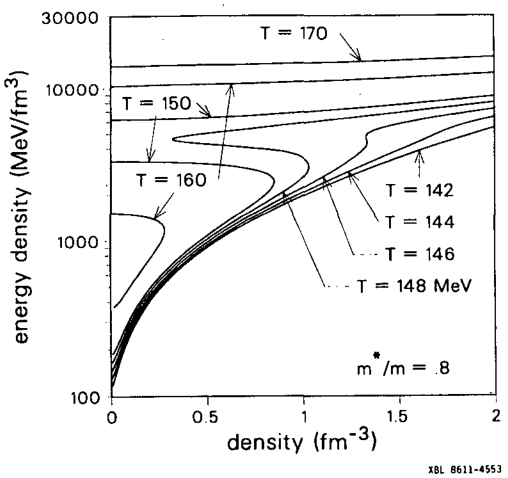

Figure 5: Isotherms of energy density as a function of net baryon density. At low temperature, only the normal state exists, while at high temperature, only the abnormal. At intermediate temperatures, around $T = 150$ MeV, both solutions exist in a finite density interval.

displace it slightly to a position of positive or negative internal pressure so as to balance the forces, but for our purpose, we can ignore this. It is metastable because its field configuration is very different from the normal state of matter, that is to say, the scalar field is very large. The baryon masses are far displaced from their vacuum values or their values in the normal state of matter. Moreover, there is a high abundance of low mass baryon-antibaryon pairs. The entropy is correspondingly very high. All of these factors introduce a barrier to its decay to the normal state. However, since the object is hot, $T \approx 150$ MeV, and since the baryons are far removed from their vacuum masses, it will radiate bosons in a spectrum characteristic of its temperature. These will include photons, lepton pairs, and pions. Since it is dense, it is plausible that the strong interactions will maintain a thermal distribution of bosons as it cools. Its spectrum will therefore be that of a very hot black-body. As it cools it will actually shrink.

The specific heat of the abnormal metastable state is obviously very high, requiring 6 GeV in radiation to cool it from 151 to 150 MeV. About two thirds of the energy of the body is emitted in black-body radiation, as it cools to $T = 147$ MeV. As it cools, baryon-antibaryon pairs, which are the glue that binds the abnormal matter, are reabsorbed, until in the vicinity of that temperature, the abnormal phase ceases to have a domain of negative pressure, and somewhat below that temperature, there is no distinction between normal and abnormal phase (see Fig. 5). The pressure then becomes positive and presumably the body will commence to expand and disassemble

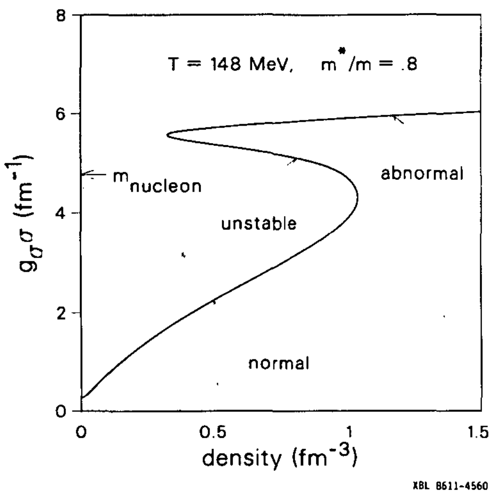

Figure 6: Scalar field strength as a function of density.

into baryons and light particle fragments. However, because of the large density separation between the two phases at the same entropy (Fig. 7), the expansion cannot be a smooth hydrodynamic one. The matter must pass through a shock discontinuity in order to conserve or increase entropy as it passes to the normal state. The energy distribution of the produced particles will not be that expected of a high energy collision because of the large fraction of the energy that was carried off in black-body radiation. Only one third of the original energy is available at the disassembly stage.

We contrast the decay of the metastable state described above with that of dense nuclear matter in the normal state or the quark-gluon plasma. For the latter two, the pressure is very high, and the disassembly will be rapid, even explosive. Much of the energy will be carried in kinetic energy of the nuclear fragments. Cooling by emission of thermal bosons will be minor. Their spectra will be strongly doppler shifted to apparent temperatures which are large compared to the actual temperatures of the expanding medium.

We expect the energy domain of the metastable state to depend sensitively on the coupling constants. These in turn are determined by bulk nuclear properties, which are well known, and on the value of the nucleon effective mass at saturation which is less well known. Therefore we examined the consequences of lowering m^*/m from 0.8 to 0.75. As expected, the domain of the abnormal state comes down appreciably in temperature and

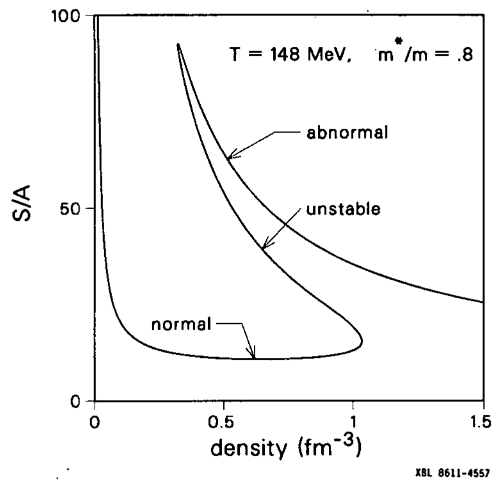

Figure 7: Entropy per baryon as a function of density.

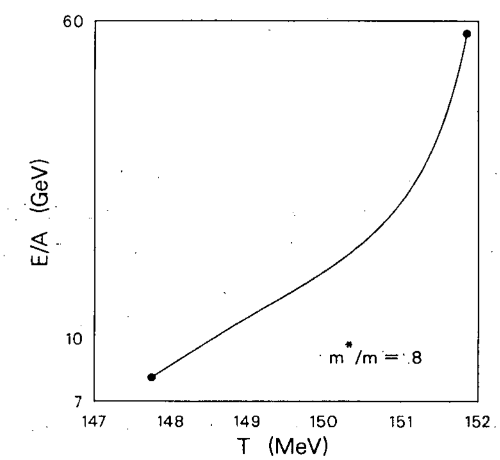

Figure 8: Center of mass energy per nucleon as a function of temperature for zero-pressure abnormal matter.

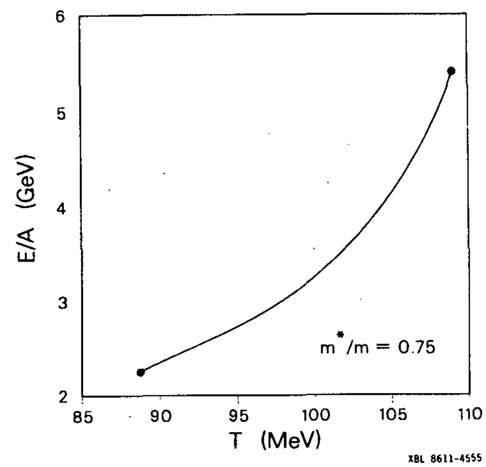

Figure 9: Center of mass energy per
nucleon as a function of temperature for zero-pressure abnormal matter for $m^*/m = 0.75$.

energy as shown in Fig. 9, with the energy range being 2 to 5 GeV. The lower value of the effective mass places the abnormal state more comfortably in the domain of validity of the theory. The discussion of this figure is similar to that above. In particular, if a fragment of such matter were produced at the upper end of its range, it would have to radiate more than half its energy in black-body spectra of bosons before it could decay into nuclear fragments. If the lower effective mass is the appropriate choice, the metastable state can be produced at much lower energy.

4 Summary

Because of their non-linearity, the field equations of nuclear field theory admit of additional solutions besides the normal state of matter. One of these is a finite-temperature abnormal phase, characterized by low baryon effective masses, abundant baryon-antibaryon pairs, and high entropy. Over a narrow range in temperature, matter can exist in the abnormal phase at zero pressure. This is a hot metastable state, for which there is a barrier against decay, because the field configuration is different than in the normal state, the baryon masses are far removed from their vacuum masses, there is an abundance of pairs also far removed from their vacuum masses, and a correspondingly high entropy. The abundance of baryon-antibaryon pairs is the glue that holds this matter together. The signals associated with this novel state are quite unusual. A fragment of such matter will cool by emitting a spectrum of black-body radiation, consisting principally of photons, lepton pairs and pions, rather than by baryon emission, because the latter are far removed from their vacuum masses. If produced at the upper end of its temperature range, a large fraction of the original energy, more than half in the examples studied here, is radiated in this way until, because of the cooling and consequent reabsorption of baryon-antibaryon pairs, which are the glue that binds the abnormal state, it reaches a domain where its pressure changes from zero to a positive value. Then it will commence to expand and disassemble, most likely through a shock discontinuity. The baryons and light elements produced in the eventual decay will account for only a fraction of the original energy.

Three aspects of the metastable state that we have not yet investigated are its cross-section for formation in nuclear collisions, its lifetime, and the increase in stability that it would acquire through the development of a net strangeness if its lifetime were long on the scale of the weak interactions. For a short-lived object produced in the collision of ordinary nuclei, the hyperons must appear in pairs. For a long lived object, some of the conserved baryon charge can be carried by hyperons, which are energetically favored when the chemical potential exceeds their effective masses. If sufficiently long-lived, an additional signal would be the appearance of hyperons in the decay products of the metastable fragment.

Concerning the prospects for the existence in nature of such a novel state as discussed here, we note that nuclear field theory is an effective one, whose range of validity is terminated at high density by the finite size of the nucleons and their underlying quark structure. Its precise range of validity is unknown. Secondly, the precise location in temperature and density of the metastable state depends sensitively on the choice of nucleon effective mass *at saturation* , which is the least well known of the nuclear properties that are used to fix the coupling constants of the theory. For $m^*/m = 0.8$, the temperature and density of the metastable state are on the marginally high side, but may still lie below the phase transition to the quark-gluon plasma. For $m^*/m = 0.75$ the metastable state lies more comfortably in the lower temperature-density domain. In any case, the theory, regarded as a mathematical model, does possess such a solution, and it would be a very novel state of matter were it to be found.

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References

[1] C.f. references in B. D. Serot and J. D. Walecka, *The Relativistic Nuclear Many-Body Problem*, in "Advances in Nuclear Physics", eds. J. W. Negele and E. Vogt, 1986;
P.-G. Reinhard, M. Rufa, J. Marhun, W. Greiner and J. Freidrich, Z. Phys. **A323** (1986) 13.

[2] B. Banerjee, N. K. Glendenning and M. Gyulassy, Nuc. Phys. **A361** (1981) 326.

[3] S. I. A. Garpman, N. K. Glendenning and Y. J. Karant, Nuc. Phys. **A322** (1979) 382.

[4] N. K. Glendenning, *Phase Transitions in Nuclear Matter* in 7th High Energy Heavy Ion Study, 1984, ed. R. Bock, H. H. Gutbrod and R. Stock (GSI report-85-10).

[5] N. K. Glendenning, Phys. Lett. **144B** (1984) 158.

[6] J. Theis, G. Grabner , G. Buchwald, J. Maruhn, W. Greiner, H. Stocker, J. Polonyi, Phys. Rev. **D28** (1983) 2286.

[7] T. Nakai and S. Takagi, Prog. Theor. Phys. **71** (1984) 1118.

[8] M. H. Johnson and E. Teller, Phys. Rev. **98** (1955) 783;

[9] H. P. Duerr, Phys. Rev. **103** (1956) 469;

[10] J. D. Walecka, Ann. of Phys. **83** (1974) 491

[11] N. K. Glendenning, B. Banerjee and M. Gyulassy, Ann. Phys. (N. Y.) **149** (1983) 1.

[12] J. Boguta and A. R. Bodmer, Nucl. Phys. **A292** (1977) 413.

[13] W. D. Myers, *Droplet Model of Atomic Nuclei* (New York: McGraw Hill, 1977);
W. D. Myers and W. Swiatecki, Ann. of Phys. **55** (1969) 395;
W. D. Myers and K-H. Schmidt, Nucl. Phys. **A410** (1983) 61.

[14] J. P. Blaizot, D. Gogny and B. Grammiticos, Nucl. Phys. **A265** (1976) 315;
J. P. Blaizot, Phys Rep. **64** (1980) 171;
J. Treiner, H. Krevine, O. Bohigas and J. Martorell, Nucl. Phys. **A317** (1981) 253.

[15] G. Co and J. Speth, Phys. Rev. Lett. **57** (1986) 547.

[16] N. K. Glendenning, Phys. Rev. Lett. **57** (1986) 1120.

[17] J. D. Walecka, Phys. Lett. **59B** (1975) 109.

[18] N. K. Glenndenning, L. P. Csernai and J. I. Kapusta, Phys. Rev. **C33** (1986) 1299.

[19] T. D. Lee and G. C. Wick, Phys. Rev. **D9** (1974) 2291.