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HIGHLY EXCITED NUCLEAR MATTER

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Introduction

In this lecture I will talk about a type of collision that will happen in only a very small fraction of the collision events produced with the Bevalac, if at all. Nevertheless, these collisions may lead to spectacular effects and be of great value in furthering our understanding of nuclear matter. The type of collision I refer to is the "head-on" collision between heavy nuclei in which there is considerable momentum exchange between the projectile and target nuclei. If the initial kinetic energy of the projectile nucleus is in the range 1-2GeV per nucleon then a large transfer of momentum between projectile and target could lead to energy densities in the center of mass system which are much larger than anything produced heretofore in the laboratory. It is this opportunity for studying matter in an entirely new regime of density and temperature that makes the search for head-on collisions very worthwhile.

If we consider the collision of identical nuclei where the incident nucleus has a laboratory kinetic energy per nucleon T then the internal energy per baryon in the center of mass system is

$$\epsilon = Mc^2 \left[\left(1 + \frac{T}{2Mc^2} \right)^{1/2} - 1 \right] \quad (1)$$

For $T = 2\text{GeV/n}$ we have $\epsilon = 400\text{MeV}$. In the case of large momentum transfer between incident and target nucleus one would expect that for a brief moment a very hot gas of baryons would be formed in which the excess center

of mass energy appears as thermal energy. If we assume that this hot baryon gas can be described as an ideal gas of nucleons then the temperature θ of this hot gas will be given by

$$\epsilon = 3\theta + \mu\epsilon^2 \left(\frac{K_1 (\mu\epsilon^2/\theta)}{K_2 (\mu\epsilon^2/\theta)} - 1 \right) \quad (2)$$

This relation differs from the familiar expression $\epsilon = \frac{3}{2}\theta$ because the nucleons must be treated relativistically. In the ultrarelativistic limit one would have $\epsilon = 3\theta$. In the neighborhood of $T = 2\text{GeV}$ eq. (2) gives $\epsilon \approx 2\theta$. The temperature as a function of T is shown in Figure 1 for the case of an idea gas of nucleons and for the case where excited states of the nucleon are taken into account (Chapline, Johnson, Teller, and Weiss, 1974).

Hot Nuclear Matter

At temperatures in excess of 100MeV mesons will be copiously produced. The rate of meson production in a proton gas as a function of temperature is shown in Fig. 2. This curve is calculated by averaging the experimental total cross-section for meson production with relativistic Maxwell-Boltzmann distributions for the colliding protons. The number of mesons produced does not, of course, keep increasing because one will soon reach a state of thermodynamic equilibrium. This state of thermodynamic equilibrium will be reached in heavy ion collisions if the decay time for the highly excited material in the center of mass system exceeds the equilibration time for meson degrees of freedom. In the "hydrodynamic" approximation the decay time of the highly excited material is approximately $.2R/c_s$ where R is the

radius of the excited material and c_s is the speed of sound in the highly excited material. The .2 factor simply accounts for the fact that half the mass lies within .2R of the surface. If we compare this decay time with the equilibration time, $(n\bar{v})^{-1}$, and assume $R = 1.1A^{1/3}$ fermi we find that the decay time exceeds the equilibration time if

$$A^{1/3} \gtrsim 4$$

Thus, collisions of nuclei with $A \gtrsim 64$ could lead to states which are in thermodynamic equilibrium with respect to meson emission and absorption.

If one assumes that the mesons are present as an ideal gas of pions then one finds that for a baryon density equal to twice normal nuclear density the equilibrium pion density is $.11(\text{fermi})^{-3}$ for $T = 2\text{GeV}/n$. This corresponds to about one pion for every three nucleons in the highly excited state. Of course, at nuclear densities above that of normal nuclei it is not a very good approximation to treat the pions present as being free because of their interactions with the nucleons. At twice normal nuclear density pions with energies inside the $\Delta(1236)$ resonance will spend about 90% of their time "resonating" with the nucleons. Thus, it might not be too bad an approximation to assume that the highly excited nuclear matter can be described as a gas of nucleons and excited nucleons.

The thermodynamic properties of a gas of nucleons and excited nucleons will be determined by the spectrum of excited nucleon states. According to the statistical bootstrap theory this spectrum has the form (Hamer and Frautschi 1972)

$$\rho(m) = \text{cm}^{-3} e^{-3 m/\theta_0} \quad (4)$$

where $\theta_0 \approx 175 \text{ MeV}$. In the region of baryon masses m where a thorough search for resonant states has been carried out, it is found that the density of states does increase rapidly with mass and could be fit with an expression of form (4). The temperature of an ideal gas of nucleons and excited nucleons, assuming that the spectrum of nucleon states has the form given in eq. (4), is shown as the solid line in Fig. 1. Because of the factor m^0 in eq. (4) the temperature of such a gas can never exceed θ_0 . That is, as was first suggested by Hagedorn (1970), nuclear matter in thermodynamic equilibrium has a maximum temperature. Since the non-equilibrium temperature (dashed line) at $T = 2 \text{ GeV}/n$ is considerably higher than θ_0 , measurement of the equilibrium temperature, perhaps by measuring the A dependence of the temperature, could provide a test of Hagedorn's hypothesis.

Shock Compression of Nuclear Matter

Since nuclei coming together at relativistic velocities are moving faster than the speed of sound in nuclear matter one would expect on the basis of classical hydrodynamics that shock waves would appear where the nuclei interpenetrated. The question of when this picture applies to collisions of nuclei is discussed in the Appendix. Here we will merely assume the hydrodynamic picture does apply and that the nuclear matter in the colliding nuclei is compressed as a result of passing through a shock wave.

The compression in a relativistic shock wave is given by (Johnson and McKee 1969):

$$\frac{n}{n_0} = \gamma \left(1 + \beta^2 \frac{e}{P} \right) \quad (5)$$

where β is the velocity of the unshocked material in the frame of the shocked material, $\gamma = (1 - \beta^2)^{-1/2}$, e is the proper energy density of the shocked material, and P is the pressure of the shocked material. In the case of identical nuclei $\gamma = (1 + T/2Mc^2)^{1/2}$ and $e = n(Mc^2 + \epsilon)$ where ϵ is given in Eq. (1).

Thus we have

$$\frac{n}{n_0} = \left(1 + \frac{T}{2Mc^2}\right)^{1/2} + \frac{nT}{2P} \quad (6)$$

for the shock compression produced by a nucleus with kinetic energy per nucleon T . The actual compression will depend on how the pressure P in compressed nuclear matter depends on density. If we neglect the interactions between baryons then $P = n\theta$ and therefore

$$\frac{n}{n_0} = \left(1 + \frac{T}{2Mc^2}\right)^{1/2} + \frac{T}{2\theta} \quad (7)$$

As an example, if the compressed nuclear matter is treated as an ideal gas of nucleons then one obtains the compressions shown in curve (b) of Fig. 3. It can be seen that one might expect to produce nucleon densities as high as $6 n_0$, i.e. $\sim 10^{39} \text{ cm}^{-3}$, with the Bevalac ($T = 2 \text{ GeV/n}$).

The actual densities produced in nuclear shock waves may be higher or lower than for an ideal nucleon gas because of the effects of meson production and nucleon-nucleon interactions. Meson production tends to lower the temperature and therefore increase the compression, as can be seen from Eq. (7). If one uses the bootstrap density of states, Eq. 4, and neglects the interaction between baryons then one obtains the compression vs. T curve shown as curve (a) in Fig. 3. It can be seen that if baryon-baryon interactions were negligible

one would expect to reach a density $n \sim 10 n_0$ at $T = 2 \text{ GeV}/n$. Of course, at densities much above normal nuclear density one does not expect that baryon-baryon interactions can be neglected. The conventional picture of nuclear forces assumes that there is a strong short range repulsion between nucleons due to ω -meson exchange. Since the ω -meson couples to the baryon number this short range repulsion would be independent of the state of internal excitation of the nucleon and therefore would lead to a lower compression than indicated by curve (a). A lower bound on the compression can be obtained by assuming that the pressure $P = ne$. The equation of state holds when the inter-nucleon separation is smaller than the range of a strong repulsive interaction (Zeldovich, 1962). It has the property that the speed of sound, i.e. $(c^2 dP/de)^{1/2}$, approaches the speed of light when the density increases to where $\epsilon \gg Mc^2$. If one substitutes $P = ne$ into Eq. (6) and uses Eq. (1) one obtains curve (c) in Fig. 3. The actual shock compression in nuclear matter will lie somewhere between curves (a) and (c) of Fig. 3.

Diagnosis of Compression

Figure 3 shows that different possible assumptions about the equation of state of compressed nuclear matter can lead to quite different values of the shock compression even at $T = 2 \text{ GeV}/n$. It is evident that a measurement of the shock compression in relativistic collisions of heavy nuclei could provide useful clues as to the nature of compressed nuclear matter. One possible indication of the compression produced in nuclear collisions would be the number of pions produced. The production of pions tends to lower the pressure and therefore increase the compression.

A rigorous statement of this effect can be arrived at by making use of the virial theorem:

$$P = n\theta - \frac{1}{6} n^2 \int_0^\infty r \frac{d\phi}{dr} g(r) 4\pi r^2 dr \quad (8)$$

where $\phi(r)$ is the nucleon-nucleon potential and $g(r)$ is the pair correlation function in nuclear matter. We are using the non-relativistic form of the virial theorem since relativistic effects will not be of great importance in the temperature regime we are considering. If we integrate the right hand side of Eq. (8) by parts we obtain

$$P = n(\theta + \epsilon_{int}) + \frac{1}{6} n^2 \int_0^\infty \phi g'(r) 4\pi r^3 dr \quad (9)$$

where $\epsilon_{int} = \frac{n}{2} \int_0^\infty \phi g(r) 4\pi r^2 dr$. The quantity ϵ_{int} is the interaction energy per nucleon due to the nucleon-nucleon potential. If we assume that the potential is positive and repulsive, then the second term of Eq. (9) will be positive and therefore, we have

$$P \geq n(\theta + \epsilon_{int}) \quad (10)$$

The equality will approximately hold when the repulsive potential is "soft", i.e. when the logarithmic derivative of $g(r)$ is small. Combining Eqs. (10) and (6) gives

$$\frac{n}{n_0} \leq \left(1 + \frac{T}{2Mc^2}\right)^{1/2} + \frac{T}{2(\theta + \epsilon_{int})} \quad (11)$$

If we assume that the repulsive potential acting between baryons is independent of the state of internal excitation, as is the case for the potential due to

ω -meson exchange, then the energy per baryon can be written

$$\epsilon = (\bar{m}(\theta) - M)c^2 + \frac{3}{2}\theta + \epsilon_{\text{int}} \quad (12)$$

where $\bar{m}$ is the average baryon mass and ϵ_{int} is independent of $\bar{m}$. We assume here that the temperature θ is low enough so that the non-relativistic approximation, $3\theta/2$, for the thermal energy can be used. If we use the bootstrap density of states, Eq. (4), to calculate $\bar{m}(\theta)$ then Eqs. (11) and (12) imply a relation between $\bar{m}(\theta)$ and n/n_0 like that shown in Fig. 4. The upper limit on the compression (solid line) holds when the repulsive potential is "soft" while the lower limit, $2\gamma + 1$, holds when the repulsive potential is "hard". It can be seen that if few pions are present in the compressed nuclear matter then the compression was near the minimum, while if many pions are present the compression can be either large or near the minimum depending on whether the repulsive potential is "soft" or "hard". Of course, as the nuclear matter expands from the compressed state it will cool leading to a reduction in the number of pions present. It is to be noted, however, that the temperature will not drop very fast if the repulsive potential is weak since the variation of temperature with volume will be determined by the ratio P/c_v where c_v is the specific heat. Since the pressure P is small for a weak soft potential the "frozen out" temperature would not differ greatly from the temperature in the compressed state. Thus measurement of the number of pions produced in the collision should give some indication whether nuclear matter is easily compressed or not.

Possibilities of Abnormal Nuclear States

In the previous discussion we have pictured compressed nuclear matter as consisting of nucleons and excited nucleons repelling each other by a repulsive potential. This is the conventional picture of compressed nuclear matter, but one should also consider the possibility that high energy collisions of nuclei could lead to abnormal states of nuclear matter. Here are a few of the possibilities:

- 1) The nucleon-nucleon potential might have an attractive region inside the repulsive core, e.g. due to f_1 exchange. This might lead to stable or metastable collapsed nuclei, i.e. systems with baryon number $A > 1$ and with radii much smaller than normal nuclei (Bodmer 1971).
- 2) If the size of excited nucleons permitted metastable nuclei containing excited nucleons might be formed. For example, one can fit 16 Δ 's into the 1S orbital without violating the Pauli principle and therefore there exists the possibility of a metastable system containing 16 Δ 's, analogous to an α -particle, but several times denser (Chapline and Weiss, 1973).
- 3) The question of whether the bootstrap density of states, eg. 3, applies to baryons in highly compressed nuclear matter is open. One possibility is that superbaryons exist, i.e. baryons with $A > 1$ and size comparable to a nucleon and for which the bootstrap density of excited states applies (Carlitz, Frantschi, and Nahn, 1973).
- 4) Metastable dense states of nuclei could exist where the binding is due to a non-vanishing scalar field (Lee and Wick, 1974).

Pion Lasers

If it turns out that excited baryon states can exist at densities higher than normal nuclear density then the rapid expansion of the compressed nuclear matter could lead to non-equilibrium states in which population inversions existed. This effect would be most pronounced in the lightest nuclei for which equilibrium with respect to excited nucleon states and also high compressions (see Appendix) can be established, since for these nuclei the drop of density with time will be most rapid. The existence of population inversions in the expanding nuclear matter will lead to stimulated emission of pions if the population inversion density N^* satisfies:

$$N^* \gg \frac{6\pi}{\lambda^2} \frac{1}{R_c} \quad (13)$$

where R_c is the radius of the compressed nuclear matter and λ is the pion Compton wavelength. That N^* must be large compared to the right hand side rather merely greater results from the fact that the pion laser would be operating without mirrors. Condition (13) suggests that stimulated emission of low energy ($\lambda < 200$ MeV/c) pions might be possible if the compressed density were 6x normal nuclear density. In fact, as Professor Teller pointed out in his talk, stimulated emission of pions may even be stronger than expected because of the possibility of double pion emission.

Applications to Astrophysics - Neutron Stars

We have seen that crude measurements of the properties of nuclear matter at densities above normal nuclear density might be made by observing head-on collisions of heavy nuclei. One area where measurements of the equation of state of dense nuclear matter should prove useful is the theory of neutron stars. Although the results on collisions of nuclei would not

be directly applicable to neutron stars because nuclei contain protons as well as neutrons, it might be possible to produce nucleon densities with the Bevalac which are comparable to the maximum nucleon densities in neutron stars. Thus experiments with the Bevalac may lead to a much better understanding of the internal structure of neutron stars. Two problems in particular which might be clarified are the maximum mass of neutron stars and the question of whether neutron stars have a solid core.

The maximum mass of neutron stars is determined by the behavior of the speed of sound in dense neutron matter (Nauenberg and Chapline, 1973). If one assumes that the equation of state of neutron matter is known up to some density n_p and that the speed of sound in neutron matter at higher densities does not exceed the speed of light, then it is possible to derive an upper limit for the mass of a neutron star (Nauenberg and Chapline loc. cit.). If the density n_p is the density of nucleons in normal nuclei then this upper limit on the maximum mass is about $3 M_\odot$. If the reference density n_p is 10^{39} cm^{-3} , i.e. ~ 6 x normal nuclear density, then the upper limit on the neutron star mass will be about $2 M_\odot$. Thus, whereas ordinary nuclear physics only allows one to infer that neutron stars cannot be more massive than about $3 M_\odot$, experiments with the Bevalac should go a long way towards defining the actual maximum mass of neutron stars, which is expected to be between 1.5 and $2.0 M_\odot$.

Experiments with the Bevalac may also clarify the question of whether neutron stars have a solid core. The solidification pressure of neutron matter depends on the "hardness" of the short range repulsive potential between nucleons. If the repulsive potential rises steeply, as it does

between atoms, then neutron matter would solidify at densities only slightly above normal nuclear density (Anderson and Palmer, 1971). On the other hand, a relatively soft potential, such as that due to vector meson exchange, would lead to solidification only at densities $\sim 10\times$ normal nuclear density (Pandharipande, 1972). Some estimates of the steepness of the rise of the repulsive potential with decreasing nucleon separation might be obtained by measuring the change in pion yield in head on collisions with energy. The "frozen out" temperatures of the pions may also be useful in this respect (see Sec. IV). It might be thought that experiments with the Bevalac would not be relevant to the question of the solidification of neutron matter because the nuclear matter in the head on collisions would be very hot. However, as a general rule experiments on shock compressed matter give information on cold compressed matter at even higher densities due to the fact the smaller interparticle separations are being sampled in the hot matter.

Applications to Astrophysics - The Early Universe

At times earlier than 10^{-4} seconds the temperature of the universe exceeded 100 MeV. At times earlier than 10^{-8} seconds the average baryon density in the universe exceeded that in nuclei. Thus while existing proton accelerators should be adequate to elucidate the physics of the universe in the period $10^{-8} \text{ sec} < t < 10^{-4} \text{ sec}$, only the Bevalac will be capable of studying directly conditions during the epoch $t < 10^{-8}$ seconds. Some of the problems that may be elucidated by experiments with the Bevalac are the origin of the cosmic black-body radiation, the origin of galaxies, and the existence of primordial black holes.

Measurements of the temperature of the cosmic black-body radiation together with estimates of the present day average baryon density in the universe yield a value for the entropy per baryon which is quite high. This suggests that strongly dissipative processes were important at some time in the early history of the universe. Since the entropy per baryon is a free parameter in cosmological models it is necessary to understand these dissipate processes in detail in order to explain the cosmic black-body radiation. It is quite possible, however, that understanding of these dissipative processes will have to await a better understanding of dense nuclear matter. For example, Zeldovich (1973) has pointed out that damping of relatively small magnitude density fluctuations in initially cold nuclear matter could account for the observed entropy if very dense nuclear matter obeys the equation of state $p = e$. Another possibility (Carlitz, Frautschi, and Nahm, 1973) is that non-equilibrium decay of superbaryons is responsible for the black-body radiation.

One particularly interesting consequence of assuming that the statistical bootstrap spectrum of excited states applies to baryons in dense nuclear matter is that large spontaneous density fluctuations would occur (Carlitz, 1972). Carlitz, Frautschi, and Nahm (1973) have suggested that these large density fluctuations in dense nuclear matter are responsible for observed structures such as galaxies and clusters of galaxies. A difficulty with this suggestion is that structure on scales larger than clusters of galaxies would tend to be washed out, which contradicts recent observations of Peebles (1974).

Another possible consequence of large density fluctuations in dense nuclear matter is that large numbers of black holes might be formed in the early universe. The existence of large numbers of small ($M \ll M_{\odot}$) black holes in the universe would resolve some current cosmological puzzles (Chapline, to be published) and might even be of economic importance. Of course, it should be kept in mind that formation of small black holes would take place at densities orders of magnitude higher than could be produced with the Bevalac. Nevertheless, any evidence from the Bevalac bearing on the question of whether the statistical bootstrap density of states applies at densities above normal nuclear densities would be valuable in assessing the plausibility of large density fluctuations in very dense nuclear matter.

Appendix

Use of the shock condition, Eq. (5), to find the compression of nuclear matter resulting from the head-on collision of heavy ions assumes that hydrodynamic equations apply. However, it is by no means obvious that the hydrodynamic approximation applies to the collision of objects containing as few particles as nuclei. In order to verify that an equation such as the shock condition, Eq. (5), applies to a given situation one must solve the Boltzmann equation. In plane geometry, one finds that it takes about 3 mean free paths to form a shock wave if the differential scattering cross-section is isotropic. For forward peaked differential cross-sections it would take more than 3 mean free paths to form a shock wave, the exact number depending on how sharply forward peaked the differential cross-section. At 2 GeV/n the differential cross-section for nucleon-nucleon scattering is so strongly peaked forward that there could be almost no hope of forming nuclear shock waves in ion-ion collisions. However, the nucleon-nucleon collisions are inelastic; the nucleons losing several hundred MeV at each collision. In the collision of two uranium nuclei, each nucleus would be ~ 7 mean free paths across so that most nucleons would be down to several hundred MeV before the two nuclei had gone completely thru each other. At several hundred MeV per nucleon lab energy the differential cross-section is nearly isotropic and therefore one would expect to get considerable momentum exchange between two uranium nuclei even at 2 GeV/n. It seems likely however, that strong momentum exchange

and the formation of shock waves is only possible in heavy nuclei, even at lower energies. Thus one can probably not use hydrodynamic equations to describe the collision of light nuclei, as was done by Scheid, Müller, and Greiner (1974). To find out what actually happens in head-on collisions of nuclei it will be necessary to solve the relativistic two-dimensional Boltzmann equation.

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Figure Captions

- Figure 1 Nuclear temperature resulting from the head-on collision of a nucleus with laboratory kinetic energy per nucleon T with an identical stationary nucleus. The dashed curve assumes that the hot nuclear matter can be described as an ideal gas of nucleons. The solid curve is for an ideal gas of N 's and N^* 's where the N^* mass spectrum has the form given Eq. 4.
- Figure 2 Thermal average rate for pion production in a hot nucleon gas, calculated using the experimental total pion production cross-section, σ , and relativistic Maxwell-Boltzmann distributions for the nucleons (Dahlbacka, Chapline, and Weaver, 1974).
- Figure 3 Nuclear shock compression that would result from the head-on collision of a heavy nucleus with laboratory kinetic energy per nucleon T with an identical stationary nucleus. Curve a (solid line) assumes that the compressed nuclear matter can be described as an ideal gas of N 's and N^* 's. Curve b (dashed line) assumes that the compressed nuclear matter can be described as an ideal gas of nucleons alone. Curve c (dotted line) assumes that the pressure of compressed nuclear matter is given by $P = e - nMc^2$.

Figure 4 Number of pions per nucleon in compressed state as a function of the shock compression. It is assumed that the compressed nuclear matter can be represented as a gas of N's and N^* 's and that N's and N^* 's all have identical interactions. The solid line corresponds to a very "soft" repulsive interaction and the dashed line corresponds to a "hard sphere" repulsive interaction. We have chosen $T = 2$ GeV/nucleon.

Figure 1

Figure 2

Figure 3

Figure 4