объединенный
институт
ядерных
исследований
дубна

E2-88-588

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EXTENDED THERMOHYDRODYNAMICS OF INTERPENETRATING NUCLEAR MATTER

Submitted to "Ядерная физика"

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1988

I. Introduction

Fluid-dynamical models have been proven to be useful means for describing averaged properties of colliding heavy nuclei at high energies. Relying on baryon and energy – momentum conservation laws and on local equilibrium concept, they visualise the average space time evolution. The equation of state and the phenomenological transport coefficients admit a large **flexibility** of such models. However, the estimate of the free mean path of nucleons /1/ point to values being comparable with the diameter of nuclei. This makes evident that the one-component hydrodynamical model in describing the high-energy collisions of heavy nuclei suffers from the a priori justification of its applicability, because it relies on the assumption of instant stopping of nuclear fluid elements just when colliding. At relativistic energies the nucleon-nucleon cross sections become foreward peaked; and at ultrarelativistic energies the longitudinal growth /2/ becomes important. **These effects** will diminish the nuclear stopping.

Already the pioneer papers (cf. /3/) on application of fluid-dynamical concepts in heavy-ion collisions, **include nuclear transparency via the two-fluid model. In the mean time this** model is modified and applied to higher energies /4/. It rests on the assumption of the two distinguishable interpenetrating nuclear fluids resembling to **both the** nuclei. One particular subtle point in the two-fluid model is the sticking process of the fluids when they are decelerated so that their relative velocity becomes smaller than the mean thermal or **Fermi** motion velocity, i.e., they are not longer distinguishable. In ref. /3/ a procedure has been introduced to go over to a one-fluid model in this situation. One possibility **of avoiding such a switching from one model to another is the application of three-fluid models** /5/.

Another approach to the posed problem is to consider the inter-penetrating nuclear matter pieces as unique system with additional internal degree of freedom, which can be visualised as the relative velocity of the nuclear subcomponents. This approach retains the usefulness of the thermo-hydrodynamical description, i.e. the notion of the nuclear equation of state whose determination is one of the ultimate goals in studying high-energy nucleus-nucleus collisions. In addition there is a smooth transition to the one fluid model.

Pioneer works in this line have been done in refs.[6,7] where anisotropic distribution functions were derived. In refs.[8] two Fermi distributions have been considered whose centres in the momentum space are separated by the relative momentum of the interpenetrating nuclei.

It is the aim of the present work to develop an extended phenomenological thermo-hydrodynamical theory which describes interpenetrating nuclear matter. The model is built up in such a way that at vanishing relative velocity (i.e. for stopped/sticked nuclei) the usual one-fluid thermo-hydrodynamics is recovered.

Some preliminary thermodynamical aspects of the present approach have been already published in refs.[9]. Here we are going to present, in a systematic way, the hydrodynamical aspects of our approach, the first steps of which have been already reported in refs.[10].

Our paper is organized as follows. In section 2 we develop the extended hydrodynamics for interpenetrating matter. In section 3 the extended thermodynamics for the Boltzmann gas is exercised. An application to colliding nuclei is discussed in section 4, where also numerical results are reported. The discussion and summary can be found in section 5. To allow a clear presentation of our basic ideas we describe the derivation of the energy-momentum tensor used in the appendices. They contain the general decomposition of the energy-momentum tensor, the specification of the entropy source density and the interpretation of certain terms entering into the energy-momentum tensor.

2. Extended hydrodynamics

For building up a closed set of dynamical equations we pose as guiding principles the following requirements:

(i) the matter is described as unique system with one additional internal degree of freedom taking into account the anisotropy of the distribution function which represents the relative velocity of the subsystems;

(ii) the distribution function is not regarded as simple sum of Maxwell Boltzmann distributions;

(iii) the two subsystems are not simply different species;

(iv) a local description applies;

(v) building elements for the dynamics are extensive densities, intensives and vector fields not higher than rank two.

We sharpen these postulates by requiring the evolution equations

$$(nu^i)_{;i} = 0, \qquad T^{ij}{}_{;j} = 0, \tag{1}$$

being the baryon conservation and energy-momentum conservation. This form used ensures the smooth transition to the usual one-fluid hydrodynamics. Eqs.(1) imply that we attribute to each fluid element such densities as baryon density n and energy density e and a four velocity u^i. According to (ii) there is no need to subdivide the elements into distinguishable subfluids (this would result in the two-fluid model). Postulate (v) excludes equations of motion of higher order tensor characteristics, such as, e.g. considered by[7]. Postulate (iii) excludes a simple diffusion approach. Since in the fluid two characteristic velocities appear, we introduce besides the time like and normalized four velocity vector u^i, a second space like vector field t^i. The latter one defines the parametrisation of the anisotropy via

$$\varsigma = n\sqrt{t^i t_i}\,. \tag{2}$$

This density is taken as new independent thermodynamical state variable; the entropy density takes therefore the form[9]

$$s = s(e, n, \varsigma)\,. \tag{3}$$

The point now is to specify the structure of the energy-momentum tensor and entropy current. In appendix A we present the general decomposition of the energy-momentum tensor. In case of a plane-symmetric motion it simplifies considerably. In appendix B the entropy source density is studied. In accordance with the second law of thermodynamics we have to require a certain structure of the time evolution law of the anisotropy density to ensure a semidefinite positive entropy production. The eigenvalues of the stress tensor are analysed in appendix C. Identifying the difference of certain eigenvectors with the anisotropy density, we arrive at an interpretation of a yet unspecified decomposition coefficient of the energy-momentum tensor. There remains one unspecified decomposition coefficient. By comparing our energy-momentum tensor with the one of two interaction-free interpenetrating Boltzmann gases (cf. appendix D), this coefficient can be interpreted as measure of the asymmetry effects. When considering the interpenetrating nuclear currents near to the symmetry plane, one can disregard the asymmetry coefficient. The energy-momentum tensor and the entropy current then read

$$T^{ij} = e u^i u^j + (p + q s_q / s_e)(g^{ij} + u^i u^j) + q t^{-2} t^i t^j, \qquad (4)$$

$$s^i = s u^i, \qquad (5)$$

and the time evolution law for the anisotropy density takes the form

$$\dot{q} = \lambda + q s_e s_q^{-1} t^{-2} t^i t^j u_{i;j}, \qquad (6)$$

$$\lambda s_q \geq 0, \qquad (7)$$

where a dot means the comoving derivative, e.g. $\dot{q} = q_{;i} u^i$. The term λ is just responsible for the entropy production. It represents the decay of the anisotropy (i.e. the "mixing" of the fluid subsystems, or in other words, the randomising of the particle momenta).

In the following we exploit the relaxation time approximation for λ,

$$\lambda = - q \tau^{-1}, \qquad (8)$$

where τ is the characteristic time for the decay of the anisotropy. Values for τ have been evaluated in ref. [7].

Concerning the proposed structure of the energy-momentum tensor some remarks are in order. (i) We restrict ourselves here onto the simplest generalising terms appropriate for the plane-symmetric motion, which is representative for the central spatial part in symmetric head-on collisions. (ii) In an approach to asymmetric collisions certainly more terms are to be included (cf. Appendix D). Also to include asymmetry effects at larger distance from the symmetry plane one needs more terms. (iii) A microscopic foundation of the energy momentum tensor is still missed. Maybe the parametrisation of the distribution functions as in refs. [8] is suited to do this. (iv) Since our generalisation does not include gradient terms, instabilities as occuring in standard dissipative relativistic fluids [11] do not appear.

3. Extended thermodynamics

The Gibbs Duhem relation

$$d(p/T) + n\, d(\mu/T) + e\, d(1/T) + q\, d(\nu/T) = 0 \qquad (9)$$

defines for the potential $s(e,n,q)$ the derived thermodynamical quantities

$$T^{-1} = s_e, \quad \mu = - s_n / s_e, \quad \nu = - s_q / s_e \qquad (10)$$

which are related according to the relation

$$p = Ts - e + \mu n + \nu q. \qquad (11)$$

The quantity T denotes the temperature, ν stands for the conjugate of the anisotropy density q, μ is the chemical potential. Eqs. (10) and (11) combine to eq. (B.5).

Now we have to specify the thermodynamical potential $s(e,n,q)$. In ref. [9] we **presented** a recipe, which goes conform with axiomatic thermodynamics, to construct, from a given potential $s_o(e,n)$, the wanted function $s(e,n,q)$ with $s(e,n,q=0) = s_o(e,n)$. Other approaches being based on certain parametrisations of the distribution function are proposed in refs. [8]. However, a concise thermodynamics is still missed for such approaches which, on the other hand, can incorporate momentum-dependent microscopic interactions. Therefore, we follow here the line of ref. [9]. To avoid too involved formulae we consider the simple case of a Boltzmann gas. Then the generalised potential $s(e,n,q)$ follows from ref. [9] as

$$s = n \left\{ \tfrac{3}{2} \log[e - m n \operatorname{ch}(q/mn)] - \tfrac{5}{2} \log n + K \right\} \qquad (12)$$

(K fixes the zero point of the entropy), while the derived thermodynamical quantities follow from eq. (10) as

$$T = \frac{2}{3} \, \frac{e - m n \operatorname{ch} \bar{\varkappa}}{n}, \qquad (13)$$

$$\frac{\mu}{T} = - \frac{3}{2} \log[e - m n \operatorname{ch} \bar{\varkappa}] + \frac{3}{2} \, \frac{m n \operatorname{ch} \bar{\varkappa} - q \operatorname{sh} \bar{\varkappa}}{e - m n \operatorname{ch} \bar{\varkappa}} + \frac{5}{2}(1 + \log n) + K, \qquad (14)$$

$$\frac{\nu}{T} = \frac{3}{2} \, \frac{n \operatorname{sh} \bar{\varkappa}}{e - m n \operatorname{ch} \bar{\varkappa}}, \qquad (15)$$

$$p = n T, \qquad (16)$$

with $\bar{\varkappa} = q/mn = \operatorname{arsh} v$, m is the nucleon mass. By construction, v is the half of the relative velocity of the nuclear subcomponents [9].

4. Dynamics of anisotropy in plane-symmetric motion

Now we are in the position to solve the dynamical equations (1). Here we want to demonstrate how the dynamics of the anisotropy and the stopping process proceeds in our approach. To avoid large-scale computer calculations we (i) restrict ourselves to the plane-symmetric motion (i.e. slab geometry) and (ii) also do not solve the momentum equation $T^{ij}_{;j}\, k_{i} = 0$ consistent with the energy conservation $T^{ij}_{;j}\, u_i = 0$ and baryon conservation $(n u^i)_{;i} = 0$; instead we use a given velocity profile and study the evolution of the anisotropy. The profile we prescribe has the form

$$u = - v_s\, x^3 \left(v_c^6 t^6 + x^6 \right)^{-1/2} \qquad (17)$$

in the rhs of the cm system. It determines in each instant of time t and at each position x a baryon velocity u (see fig.1). Notice that, according to $n u^i = n_1 u_1^i + n_2 u_2^i$, the vanishing of the common velocity field u does not mean the vanishing of the velocities $u_{1,2}$ of the nuclear subsystems.

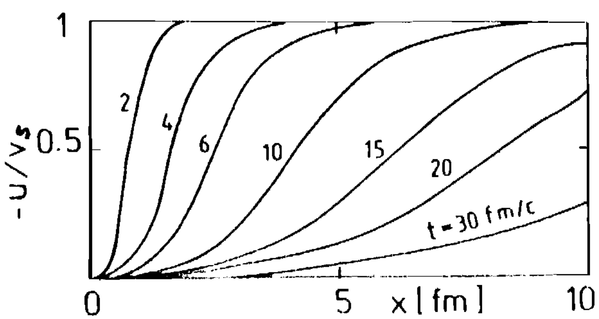

Fig.1. The velocity u (being the spatial component of the four velocity) in units of v_s. The constant v_c (cf. eq.(17)) is chosen as 0.5.

We exploit comoving synchronised coordinates

$$\dot{x}(\xi, t) = u(x, t) / \gamma(x, t) \qquad (18)$$

which define the trajectory of a fixed fluid element with initial position $x(\xi, t=0) = \xi$ via $x(\xi, t)$. The solution of eq. (18) is displayed in fig.2.

In the present coordinate system the remaining dynamical equations take the form

$$\dot{n} = -n\, \Theta / \gamma \,, \qquad (19)$$

$$\dot{e} = -\left(e + p + q\left[1 + s_q / s_e \right] \right) \Theta / \gamma \,, \qquad (20)$$

$$\dot{q} = -q \left(\tau^{-1} + \Theta / v \right) / \gamma \,, \qquad (21)$$

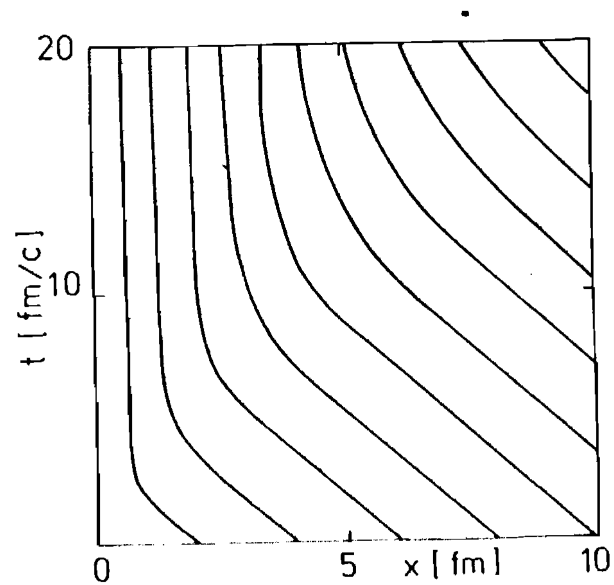

Fig.2. The trajectories $x(\xi, t)$ for the given velocity profile displayed in fig.1. Note that the trajectories belong to the total baryon current and not to that of a nuclear subsystem.

with $\Theta = \left(\frac{u}{\gamma} \partial_t + \partial_x \right) u$ as expansion; $\gamma = \sqrt{1 + u^2}$. In comparison with the standard hydrodynamics, in eq. (20) the enthapy is modified. The evolution of q is determined by the driving term $q\,\Theta$ $(\Theta < 0)$ accounting for the interpenetration in building up the anisotropy, and the term $q\,\tau^{-1}$ accounting for the decay. We assume here a relaxation of the anisotropy via randomisation in binary collisions. The relaxation time has been calculated in ref. /7/ as $\tau = (n\,\eta)^{-1}$ and $\eta = 2\,\mathrm{fm}^2$ as representative value. Solutions of the eqs. (18)–(21) are displayed in figs. 3–5. One observes indeed the expected building up of the anisotropy q and afterwards its decay. Parallel going is the half of the relative velocity of the nuclear subsystems v. In the given example with $v_s = 0.7$ (corresponding to $E_{lab} \approx 2$ GeV A) one finds typically times of 10 fm/c up to the vanishing of the anisotropy being indicative for stopping, i.e. $v \to 0$ (cf. fig.3). Quite surprising is the fact that the relative velocity remains so much below the incoming velocity v_s. This is a consequence

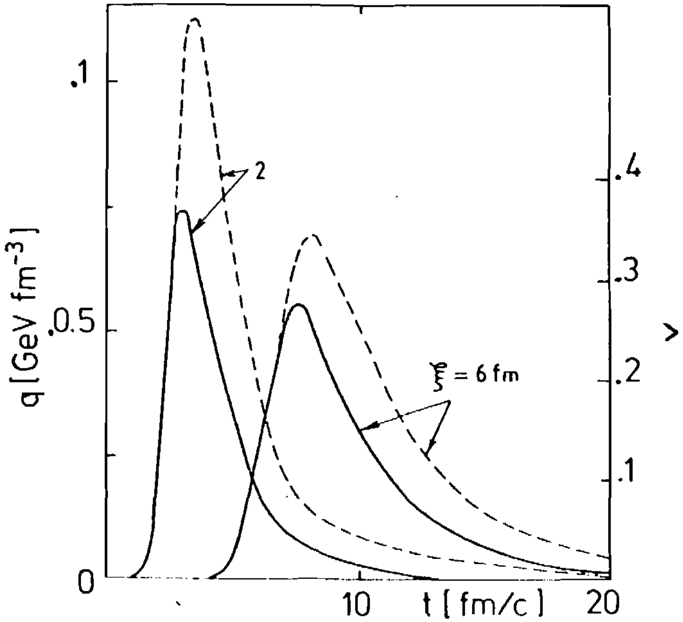

Fig.3. The time evolution of the anisotropy ξ (dashed lines) and the half of the relative velocity of the **interpenetrating** nuclei v (full lines) for the comoving fluid elements with initial position ξ = 2 and 6 fm.

that from the very **beginning, the** relaxation of anisotropy prevents the system to develop a large v. In the limit $\tau \rightarrow \infty$ we recover indeed $v \rightarrow v_s$.

The only dissipative mechanism included in our model is the relaxation of anisotropy. It is responsible for the entropy increase and, therefore, the temperature rises smoothly (see fig.4). If stopping is achieved, then the temperature remains constant. Conform to stopping and temperature rising the density is increased too. It should be noticed that due to the prescription of the velocity profile and the treatment of semiinfinite slabs the numerical values of T and n should not be taken too literally. Accounting for finite nuclei and momentum conservation will diminish these values. Also we expect that asymmetry effects play an important rôle already at a distance of a few fm beside the symmetry plane $x = 0$. (This can be seen from the snapshots of the density, temperature and relative velocity profiles in Fig.5).

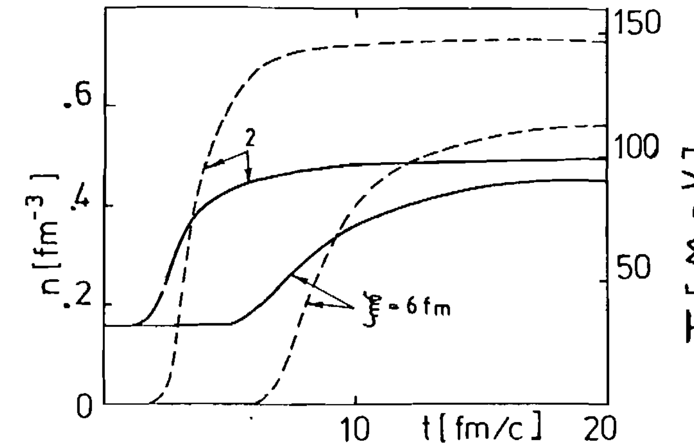

Fig.4. The time evolution of the baryon density n (full lines) and the temperature T (dashed lines) in the comoving fluid elements being initially at ξ = 2 and 6 fm.

Having solved the dynamical problem, one can reconstruct the motion of the nuclear subsystems. Due to our symmetry assumption, $n_1 = n_2$ holds; the velocities of the subsystems in the coordinate system, where the total baryon current vanishes, are, by definition, $\pm v(1+v^2)^{-1/2}$. Transforming into the centre of mass system, where the velocity profile u(x,t) is given, we get for the velocities of the subsystems

$$\bar{u}_{1,2} = \frac{\bar{u} \pm v(1+v^2)^{-1/2}}{1 \pm \bar{u}\, v(1+v^2)^{-1/2}}, \qquad (22)$$

where $\bar{u} = u(1+u^2)^{-1/2}$. Solving similar equations as eq. (18), $\dot{x}_{1,2} = \bar{u}_{1,2}$, one gets the trajectories $x_{1,2}(\xi, t)$ of the subsystems. It can be seen, however, from fig.5 that, even for rather large values of time, the profiles are far from being space independent. Therefore, the asymmetry term is not longer negligible to obtain a realistic space time picture of the collision process. Thus we do not **intend** to reconstruct the **trajectories** of the nuclear subsystems.

5. Summary

In the present paper we develop an approach to a relativistic phenomenologically extended hydrodynamics which is aimed to describe the interpenetration of nuclei. In this way we have regard to the

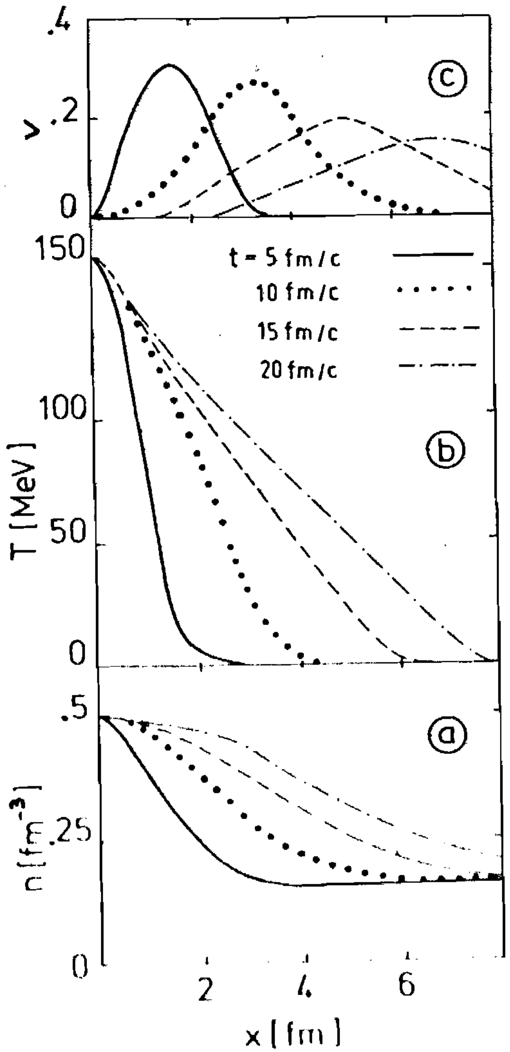

Fig.5. Snapshots of the density profile (a), the temperature (b) and the half of the relative velocity (c) at different times.

(at least partial) transparency of nuclei which stems from the finite mean free path of nucleons which is not negligible in comparison with the nuclear diameter.

The matter is regarded as unique thermodynamical system with additional internal degree of freedom. The latter one takes into account the relative velocity of the interpenetrating subsystems. The stopping process proceeds smoothly via the gradual decrease of the relative velocity. In a microscopic approach this relative velocity can be ascribed to the anisotropy of the distribution functions; the corresponding relaxation time for its decay is then accessible, e.g. in exploiting the Boltzmann equation.

The main point is the structure of the energy-momentum tensor. The general structure of it is presented. In our first explorative study we have neglected asymmetry terms. A simplified numerical illustration example shows that for a realistic description of symmetric head-on collisions the asymmetry term must be included to obtain a satisfying picture of the space time evolution. Relying on previous estimates of the relaxation time, we find typical time scales of 10 fm/c for the stopping of interpenetrating nuclear matter streams. Thus only heavy nuclei have a good chance to be stopped and to approach to local equilibrium during their collision process.

Acknowledgement

The authors are indebted to Katalin Martinas for helpful discussions concerning the thermodynamical aspects of our work. Stimulating discussions with V.D. Toneev and I.N. Mishustin are gratefully acknowledged.

Appendix A:

The decomposition of the energy-momentum tensor

The general decomposition of T^{ij} reads

$$T^{ij} = \alpha u^i u^j + \beta (u^i t^j + t^i u^j) + \gamma t^i t^j + a^i u^j + u^i a^j + b^i t^j + t^i b^j + k k^{ij} + c^{ij}, \tag{A.1}$$

$$a^i u_i = a^i t_i = b^i u_i = b^i t_i = 0,$$

$$c^{ij} u_i = c^{ij} t_j = 0.$$

The quantities b^i and d^i are yet unspecified expansion vectors. The projector

$$h^{ij} = g^{ij} + [t^2 u^i u^j + t^i t^j - (u^l t_l)(u^i t^j + t^i u^j)] \cdot (t^2 + (u^l t_l)^2)^{-1} \tag{A.2}$$

fulfills as usual $h^{ij} h_{jl} = h^i_l$ and $h^{ij} u_j = h^{ij} t_j = 0$. The entropy current is accordingly decomposed

$$s^i = s u^i + a t^i + b b^i + d d^i. \tag{A.3}$$

To identify the many terms in eq. (A.1) we firstly notice that the energy conservation $T^{ij}{}_{;j} u_i = 0$ leads to the balance equation in the form $(e u^i)_{;i}$ = source terms, i.e.

$$(e u^i)_{;i} = (\beta t^i + d^i + c^{ij} u_j)_{;i} - u_{i;j} (\hat{\gamma} t^i t^j + \beta t^i u^j + d^i u^j + b^i t^j + u^i d^j + t^i b^j + c^{ij} + h h^{ij}), \tag{A.4}$$

$$e = \alpha - \beta u^l t_l .$$

Without microscopic support and/or serious simplifications the further interpretation of terms in eq.(A.1) is not possible. Therefore, we specialize to the case of slab geometry often used as first order approximation for head-on collisions. This implies (i) $b^i = d^i = 0$ since there are no gradients or flows into the perpendicular direction; (ii) $c^{ij} = 0$ due to the isotropy in perpendicular direction. In addition we choose the particular "gauge" of the yet unspecified vector t^i in such a way that $u^i t_i = 0$.

Appendix B: The entropy source density

The second law of thermodynamics requires

$$s^i{}_{;i} \geq 0. \tag{B.1}$$

Exploiting eqs. (A.3) with (A.4),(A.5) and the baryon conservation (1) and the form of the entropy density eq. (3), we find

$$s_q \dot{q} + (s - s_e e - s_n n - s_e h) + t^i (s_e \beta_{,i} + \alpha_{,i})$$

$$+ t^i{}_{;i} (\alpha + s_e \beta) + u_{i;j} t^i t^j s_e (h t^{-2} - \hat{\gamma})$$

$$- u_{i;j} u^i t^i s_e \beta \geq 0. \tag{B.2}$$

To ensure a semidefinite positive entropy production we require for the time evolution of the anisotropy density $q = n t$, $t = (t^i t_i)^{1/2}$,

$$\dot{q} = \lambda + \nu (T_{,i} + T u_{i;j} u^j) t^i + \vartheta u_{i;j} t^i t^j, \tag{B.3}$$

where we use $T = s_e^{-1}$ as shorthand notation (in standard thermodynamics T is the temperature). (It is worth to notice that the introduction of a time evolution law for the anisotropy parameter t instead of the anisotropy density q, would result in slightly different formulae). Due to the independence of the differently contracted vectors in eq. (B.2) this equation is fulfilled if

$$s_q \lambda \geq 0,$$

$$\vartheta = s_e (\hat{\gamma} - h t^{-2}) / s_q, \tag{B.4}$$

$$\alpha = - s_e \beta = - s_q \nu / s_e ,$$

$$h = p + q s_q / s_e ,$$

where the thermodynamics pressure obeys

$$s - s_n n - s_e e - s_q q - s_e p = 0. \tag{B.5}$$

Appendix C. The coefficient ϑ

Consider the stress tensor

$$p^{ij} = p (g^{ij} - u^i u^j) + \vartheta s_q s_e^{-1} t^i t^j. \tag{C.1}$$

In general it possesses four eigenvectors which are orthogonal. Two of them are

(i) u^i with eigenvalue 0,

(ii) t^i with eigenvalue $s_q s_e^{-1} t^2 + p$.

The remaining eigenvectors are both orthogonal to u^i and t^i, therefore the eigenvalues are degenerate with the value p. Now we identify the anisotropy density q with the difference of the eigenvalues corresponding to eigenvectors in t^i direction and u^i, t^i direction, i.e.

$$q = \vartheta s_q s_e^{-1} t^2 \tag{C.2}$$

thus having arrived at an interpretation of ϑ . The energy-momentum tensor now reads

$$T^{ij} = e u^i u^j + (p + q s_q / s_e)(g^{ij} + u^i u^j) + \nu s_q s_e^{-2} (u^i t^j + t^i u^j) + q t^{-2} t^i t^j \tag{C.3}$$

and the entropy current is

$$s^i = s u^i - \nu s_q s_e^{-1} t^i . \tag{C.4}$$

Appendix D. The coefficient β

To interpret the coefficient β (and also α and ν) let us consider two non-interacting interpenetrating Boltzmann gases. Then the first and second momenta of the distribution function define the baryon current and energy-momentum tensor as

$$n^i = n u^i = n_1 u_1^i + n_2 u_2^i , \tag{D.1}$$

$$T^{ij} = w_1 u_1^i u_1^j + w_2 u_2^i u_2^j + (p_1 + p_2) g^{ij}, \tag{D.2}$$

where $w = e + p$, and the indices 1,2 belong to the two components. Eq. (D.1) relates the velocities via

$$u^i = (n_1 u_1^i + n_2 u_2^i)/n , \tag{D.3}$$

$$n = (n_1^2 + n_2^2 + 2 n_1 n_2 u_1^i u_{2i})^{1/2} . \tag{D.4}$$

Now we introduce the vector field t^i

$$t^i = \hat{\vartheta} (u_1^i + u_2^i \hat{\tau}) , \tag{D.5}$$

where

$$\hat{\tau} = - \frac{n_1 + n_2 u_1^i u_{2i}}{n_2 + n_1 u_1^i u_{2i}} \tag{D.6}$$

ensures the "gauge" $u^i t_i = 0$. The length of t^i is parametrized by $\hat{\vartheta}$ which does not concern us. Substituting $u_{1,2}^i$ by u^i and t^i we arrive at

$$\begin{aligned} T^{ij} &= \bar{\alpha} u^i u^j + \bar{\beta} (u^i t^j + t^i u^j) + \bar{\gamma} t^i t^j + \bar{k} (g^{ij} + u^i u^j - t^{-2} t^i t^j), \\ \bar{\alpha} &= n^2 N^{-2} (\hat{\tau}^2 w_1 + w_2) - \bar{k} , \\ \bar{\beta} &= - n N^{-2} \hat{\vartheta}^{-1} [\hat{\tau} n_2 w_1 + n_1 w_2], \\ \bar{\gamma} &= N^{-2} \hat{\vartheta}^{-2} [n_2^2 w_1 + n_1^2 w_2] + \bar{k} t^{-2} , \\ \bar{k} &= p_1 + p_2 , \quad N = n_2 - n_1 \hat{\tau} . \end{aligned} \tag{D.7}$$

In the symmetric case, $n_1 = n_2$, $w_1 = w_2$, $p_1 = p_2$, we have $\hat{\tau} = -1$ and accordingly $\bar{\beta} = 0$. This shows that the coefficient $\bar{\beta}$ is a measure of the asymmetry effects. For the central part in symmetric collisions, $\beta = \nu = \alpha = 0$ is therefore a suitable approximation. In this approximation we also arrive at an interpretation of the vector t^i as difference of the four velocities (up to a normalisation factor, cf. eq. (D.5)).

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Received by Publishing Department
on August 1, 1988.

Кэмпфер Б. и др. E2-88-588
Обобщенная гиродинамика взаимопроникающих потоков ядерной материи

Представлено феноменологическое обобщение релятивистской термогидродинамической модели, учитывающее взаимное проникновение ядер и, тем самым, конечность среднего свободного пробега нуклонов. Тормозная способность ядер при их высокоэнергетическом столкновении проявляется как непрерывный процесс. Элементы объема, содержащие ядерное вещество из обеих взаимопроникающих подсистем, рассматриваются как единая материя с дополнительной внутренней степенью свободы, характеризуемой относительной скоростью подсистем или анизотропией функции распределения. Динамика этой дополнительной степени свободы учтена соответствующим образом. Показано, как непрерывный характер тормозной способности ядер проявляется в данном подходе.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1988

Kämpfer B. et al. E2-88-588
Extended Thermohydrodynamics of Interpenetrating Nuclear Matter

A relativistic phenomenologically extended thermohydrodynamical model is presented which allows for the mutual interpenetration of nuclei, thus taking into account the finiteness of the nucleon mean free path. The stopping of nuclei during their high energy collision process appears as gradual process. Fluid elements with two interpenetrating subsystems are treated as unique matter with an additional internal degree of freedom, being the representative of the relative velocity of the subsystems or the anisotropy of the distribution function. The dynamics of this additional degree of freedom is properly included. We demonstrate how the smooth stopping of nuclei is described in the present approach.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna 1988