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The Phase Transition to the Quark–Gluon Plasma and Its Effect on Hydrodynamic Flow[†]

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Abstract

It is shown that in ideal relativistic hydrodynamics a phase transition from hadron to quark and gluon degrees of freedom in the nuclear matter equation of state leads to a minimum in the excitation function of the transverse collective flow.

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Hydrodynamics represents (local) energy, momentum, and charge conservation [1, 2, 3]. Because of its simplicity it has found widespread application in studying the dynamical evolution of heavy–ion collisions (see e.g. [2, 3, 4] and refs. therein). For instance, 3+1–dimensional hydrodynamical calculations for collisions at BEVALAC energies ($E_{Lab}^{kin} = 0.1 - 2$ AGeV) were performed about twenty years ago (see the very detailed review of this topic in [2]). It was found that the compressional shock waves created in the collision lead to collective flow phenomena like sideward deflection of matter in the reaction plane ("side-splash" and "bounce-off") as well as azimuthal deflection out of the reaction plane ("squeeze-out"). The confirmation of these collective flow effects by BEVALAC experiments [5, 6] was one of the main successes of the fluid-dynamical picture.

Inspired by this success, 3+1–dimensional ideal relativistic hydrodynamics was also applied to study heavy–ion collisions at AGS ($E_{Lab}^{kin} \simeq 10 - 15$ AGeV) and even CERN–SPS energies ($E_{Lab}^{kin} \simeq 60 - 200$ AGeV) [4, 7, 8, 9]. Since ideal hydrodynamics assumes that matter is in local equilibrium at every instant, colliding fluid elements are forced by momentum conservation to instantaneously stop and by energy conservation to convert all their kinetic energy into internal energy (compression and heating via shock waves). The longitudinal rapidity loss in individual nucleon–nucleon collisions is, however, limited. Thus, immediate complete stopping is not achieved in reality and, for higher beam energies, it is no longer justified to treat the initial stage of the reaction in an ideal hydrodynamical picture. Ideal hydrodynamics might nevertheless be applicable in the expansion stage of the collision [10], where the conditions of local thermodynamical equilibrium are more likely to be established. In order to describe the initial stage, however, one has to account for non-equilibrium effects [3].

These effects are naturally accounted for in microscopic transport models like QMD [11], RQMD [12], ARC [13], or ART [14]. For BEVALAC, SPS, and AGS energies, these models explain almost all single particle observables satisfactorily in terms of hadronic physics (see, for instance, [11, 15] and refs. therein). However, the basic assumption of these models is that the original (quantum) many–body problem can be adequately decomposed in terms of (classical) two–particle scatterings. For heavy–ion collisions, this assumption is up to now unproven. Moreover, to describe these two–particle scattering events the measured free hadron–hadron cross sections are employed, at least as far as these are known experimentally. This leaves a large uncertainty with respect to unknown cross sections for heavier mass resonances and effects of the nuclear environment. Given the fact that microscopic models deal with this uncertainty by introducing a large number of parameters, it is highly desirable and of considerable interest and importance to investigate, to what extent the simpler (and well defined) hydrodynamical theory is applicable to describe heavy–ion collisions in the beam energy range from 0.1 to 10 AGeV.

The simplicity of the (ideal) hydrodynamical approach lies in the fact that the only physical input is the nuclear matter equation of state (EoS) which is calculable by means of thermodynamics [2]. This constitutes also a main advantage of hydrodynamics over microscopic models, since e.g. the phase transition to the quark–gluon plasma (QGP) [16] can be studied in a simple, straightforward manner[1], whereas microscopic transport models require an ad hoc deconfinement–hadronization mechanism [18]. As we shall see, the phase transition to the QGP may indeed play a decisive role for the dynamical evolution of the system in the beam energy range $E_{Lab}^{kin} = 0.1 - 10$ AGeV.

As discussed in [19, 20] (see also [21, 22]), a phase transition in the EoS qualitatively changes the hydrodynamical flow pattern. In particular, the hydrodynamically stable solution for one-dimensional, stationary compression of matter is no longer a single shock wave, but a sequence of shock and compressional simple waves. Analogously, the hydrodynamically stable solution for one-dimensional, stationary expansion is a sequence of simple rarefaction waves and rarefaction shock discontinuities, instead of a single simple rarefaction wave.

In [23] it is shown (for matter without a conserved charge) that these phenomena occur not only if the EoS has a first order phase transition, but also if the energy (or entropy) density rises sufficiently rapidly as a function of the temperature. But even if this increase is only moderate, and a simple rarefaction wave is the hydrodynamically stable expansion solution, the flow pattern shows structures resembling broadened ("smeared") versions of rarefaction shocks. Decisive for their occurrence is the existence of a so-called "softest point" in the EoS [24], i.e., a local minimum of p/ϵ as a function of ϵ (p is the pressure, ϵ the energy density in the local rest frame of matter).

As pointed out in [19, 24], the existence of this "softest point" leads to a prolonged expansion of matter and consequently to a long lifetime of a mixed phase of QGP and hadron matter. Analogously, it also takes longer to compress matter in the early stage of a heavy–ion collision [20]. In this work we shall demonstrate in detail how both effects lead to a *minimum in the excitation function of the directed transverse collective flow in heavy–ion collisions.* An observation of this phenomenon would be a clear signature for a change of nuclear matter properties as for instance in the transition from hadron to quark and gluon degrees of freedom. We remark that our results are in agreement with earlier works [8, 9] (see also [25]) which indicated a decrease of the transverse collective flow if the nuclear matter EoS features a phase transition to the QGP.

The relativistic hydrodynamical equations represent (local) energy–momentum conservation

$$\partial_\mu T^{\mu\nu} = 0 \tag{1}$$

and (local) charge conservation

$$\partial_\mu N^\mu = 0 \ . \tag{2}$$

Here $T^{\mu\nu}$ is the energy–momentum tensor and N^μ the net baryon number current. In case that also other conserved quantum numbers are to be considered (like e.g. net strangeness) there is an additional equation of the type (2) for each of these charges. Provided that matter is in local thermodynamical equilibrium, the energy–momentum tensor $T^{\mu\nu}$ and the baryon current N^μ assume ideal fluid form [26], i.e.,

$$T^{\mu\nu} = (\epsilon + p) u^\mu u^\nu - p g^{\mu\nu} \ , \tag{3}$$

$$N^\mu = n u^\mu \ , \tag{4}$$

[1]Non-equilibrium phenomena like supercooling and bubble formation [17] are neglected in this picture.

where n is the baryon number density in the local rest frame of the fluid, $u^\mu \equiv \gamma(1, \mathbf{v})$ is the fluid 4–velocity ($\gamma \equiv (1-\mathbf{v}^2)^{-1/2}$, $\mathbf{v}$ is the fluid 3–velocity), and $g^{\mu\nu} = \mathrm{diag}(+,-,-,-)$ is the metric tensor. The equations of ideal fluid-dynamics are closed by specifying the nuclear matter EoS in the form $p = p(\epsilon, n)$.

This EoS is constructed as follows (for details see [20]). For the QGP phase, the MIT bag EoS [27] (for massless gluons and u and d quarks) is employed (with a Bag constant $B = (235\,\mathrm{MeV})^4$), while the hadronic phase is described by a version of the $\sigma-\omega$–model [28] (plus massive thermal pions) which features more realistic values for the ground state incompressibility ($K_0 \simeq 300$ MeV) and the effective nucleon mass ($M_0^* \simeq 0.635\,M$, where M is the free nucleon mass) than the original version proposed by Walecka (where $K_0 \simeq 550$ MeV, $M_0^* \simeq 0.54\,M$) [29]. Both equations of state are matched via Gibbs' conditions of phase equilibrium. Thus, the resulting nuclear matter EoS has by construction a first order phase transition between hadron and quark–gluon matter. In Fig. 1 we show the pressure as a function of energy density and baryon number density. In the QGP phase the pressure is independent of n, since for the bag model one has the simple relationship $p = (\epsilon - 4B)/3$. The mixed phase is distinct from the other phases in that the pressure is only slowly varying with ϵ and n. The hadronic phase corresponds to the small strip between mixed phase and unphysical (ϵ, n)–combinations where the pressure is zero[2]. In the following we will compare results obtained with this EoS to those obtained with an EoS for *pure* hadronic matter, i.e., matter described *solely* by the above mentioned hadron matter EoS.

Let us first investigate compression of nuclear matter in a simple *one-dimensional* "slab-on-slab" collision. The hydrodynamical solution to this problem was discussed in detail in [20]. In Fig. 2 we present various thermodynamic quantities in the final compressed state as functions of the beam energy (of a fixed–target experiment) for the above described EoS with phase transition (full line) in comparison to quantities obtained in a single shock compression of pure hadronic matter (dashed line). One clearly observes in Fig. 2(c) that in the case of a phase transition the pressure does not increase as fast with the beam energy as in the pure hadron matter case. The reason is that in the mixed phase the pressure increases only slowly as a function of energy density and baryon number density, cf. Fig. 1. This leads to the existence of a "softest point" at the phase boundary between mixed and quark–gluon matter, cf. Fig. 2(e), corresponding to $E_{Lab}^{kin} \simeq 4.1$ AGeV. As a consequence, for the same beam energy matter is much easier to compress, i.e., larger energy and baryon number densities can be obtained than for the "stiff" pure hadron matter EoS, cf. Figs. 2(a,b). Note that the specific entropy $\sigma \equiv s/n$ in Fig. 2(d) is (approximately) constant as a function of the beam energy in the mixed phase. Assuming subsequent adiabatic ($\sigma = const.$) expansion of matter until freeze-out, it was shown in [30] that a corresponding plateau should occur in the excitation function of the pion multiplicity. This plateau could serve as a signature for QGP formation.

We remark that for the above constructed EoS the phase transition sets in at $E_{Lab}^{kin} \simeq 1.5$ AGeV and pure QGP is formed at $E_{Lab}^{kin} \simeq 4.1$ AGeV. These values appear rather low. The reason is that the hadronic part of our EoS features only nucleons and pions. As a conse-

[2]Note that an ideal gas of nucleons has a minimum energy density for a given baryon density, namely the Fermi energy density at $T = 0$.

quence, the pressure as a function of temperature and chemical potential is still quite small at the phase transition. If more hadronic resonances are included, to first approximation (i.e., assuming them to behave as ideal gases) their partial pressures add to the total pressure. This shifts the phase transition to larger values of temperature, chemical potential, and consequently to higher energy and baryon number densities. This will in turn also shift the onset of the phase transition to higher beam energies (~ 10 AGeV according to the results of [31]). However, apart from such a shift the results of Fig. 2 should remain qualitatively unchanged.

Fig. 3 shows the time (in the CM frame of two equal nuclei) the compression waves need to completely compress the incoming nuclei [20]. One observes that this time is prolonged if the EoS features a phase transition (full line) as compared to the case where pure hadronic matter is compressed (dotted line). The reason is intuitively clear, since a higher compression reduces the velocity of the compression fronts travelling into uncompressed matter. The higher compression will also cause the compressed system to occupy a smaller spatial volume.

Let us now study the expansion stage in the one-dimensional scenario. For the qualitative arguments to be presented, it is sufficient to restrict considerations to baryon-free matter and, since in this case hadronic matter consists predominantly of (thermal) pions, to a massless pion gas for the hadronic part of the EoS. This is a good approximation for temperatures below 200 MeV. Above that value, nucleons become massless in the $\sigma-\omega$–model [32] and contribute roughly the same amount to thermodynamic quantities as pions. For our choice of the Bag constant, however, the phase transition to the QGP happens prior to this phenomenon ($T_c \simeq 170$ MeV).

Employing this EoS for baryon-free matter, we solve the hydrodynamic equations (1, 2) with the relativistic HLLE algorithm [19]. The initial condition is a blob of size $2\,R$ and constant energy density ϵ_0. In Fig. 4 we show the lifetime (in the CM frame) of the mixed phase in the center of the compressed system as a function of ϵ_0 [19]. One observes a pronounced peak around the "softest point" of the EoS (near the energy density ϵ_Q, i.e., at the boundary between mixed phase matter and the QGP). How this peak is related to the hydrodynamic expansion solution was explained in detail in Ref. [19]. The reason for this prolongation of the lifetime is that the system does not expand (and cool) as rapidly as in the case without a phase transition, but stays in a comparatively small spatial volume for a long time. The rapid cooling for the case without phase transition is illustrated by the dashed line which shows the CM time when the temperature of the *pure* pion gas drops below the temperature $T_c \simeq 170$ MeV in the center of the system[3].

To conclude our investigation of the simple one-dimensional scenario, the presence of a transition from hadron to quark and gluon degrees of freedom in the nuclear matter EoS leads (a) to a prolonged compression stage where the final values for energy and baryon number density are larger than for the case without phase transition. Consequently, the zone of compressed matter occupies a smaller spatial volume. It furthermore leads to (b) a prolonged expansion stage, i.e., matter will not expand and cool rapidly but stay in that relatively small spatial volume for a long time.

[3]The notion of a "phase transition temperature" is, of course, irrelevant in this case.

Let us now turn to 3+1–dimensional hydrodynamical calculations. We first consider a Au+Au–collision at 5 AGeV (i.e., close to the "softest" point in the EoS, cf. Fig. 2(e)) and finite impact parameter $b = 3$ fm. For all multi-dimensional calculations presented here we use a SHASTA algorithm with first order accuracy in time [19, 33]. The grid spacing is chosen as $\Delta x = 0.3$ fm, the time step width is $\Delta t = 0.4\,\Delta x$. The use of this rather coarse grid spacing [19, 20] and the first order scheme (which is not as accurate as a second order scheme [19]) is a concession to the enormous calculational effort of 3+1–dimensional hydrodynamical calculations. For the results presented below, however, we do not expect major quantitative changes when using a finer grid and a second order scheme.

In Fig. 5 we show CM frame baryon density contours (and flow velocity vectors) in the reaction plane at different CM times. Part (a) employs the EoS with phase transition and part (b) the pure hadronic EoS. One notices the following distinct feature: for the EoS with phase transition the compressed zone in the center of the reaction grows and expands much more slowly (and also the compression is much higher) than in the collision calculated with the pure hadronic EoS. Since this is completely analogous to our findings in the one-dimensional scenario, the previous considerations immediately present the obvious explanation for this phenomenon also in three space dimensions. In the latter case, however, this has the following further consequence for spectator matter: since compressed matter does not expand rapidly in the case of the EoS with phase transition, it does not exert pressure onto the spectators, which consequently pass the participants undeflected. In contrast, for a pure hadronic EoS the compressed zone expands violently and deflects spectator matter (the mentioned "bounce–off" effect [2]).

This effect can be seen best in the mean transverse (in–reaction–plane) fluid momentum per baryon as a function of longitudinal fluid rapidity,

$$\langle p_x/N\rangle(y) \equiv \frac{\sum_i' N_i \; p_{x,i}}{\sum_i' N_i} \; . \tag{5}$$

Here, the sum is over all fluid elements i (i.e., in praxi numerical grid cells) subject to the constraint that their longitudinal rapidity $y_i \equiv \frac{1}{2}\ln[(1+v_{z,i})/(1-v_{z,i})]$ obeys $y - \Delta y/2 \leq y_i \leq y + \Delta y/2$ ($v_{z,i}$ is the z–component of the 3–velocity of fluid element i), and N_i is the total net baryon number contained in fluid element i. Furthermore, $p_{x,i} \equiv M\gamma_i v_{x,i}$ (here $\gamma_i = (1-\mathbf{v}_i^2)^{-1/2}$). We remark that (5) is not directly comparable to an experimentally measurable quantity, since $\langle p_x/N\rangle(y)$ is essentially (a constant factor, M, times) the average x–component of the *fluid* 4–velocity in the rapidity bin $[y-\Delta y/2, y+\Delta y/2]$, and *not* the x–component of the average nucleon momentum. The latter can only be obtained after a suitable freeze–out procedure is applied to the hydrodynamic quantities. This is the subject of a forthcoming paper [34].

In Fig. 6 we show the quantity (5) at the end of the collision (defined as the CM time when the baryon number (fluid) rapidity distribution does no longer change appreciably) as a function of rapidity y normalized to beam rapidity y_{CM} for Au+Au–collisions at $b = 3$ fm and for beam energies (a) $E_{Lab}^{kin} = 3.5$ AGeV, (b) 5 AGeV, and (c) 11.7 AGeV. One observes that for the pure hadronic EoS (open circles) the $\langle p_x/N\rangle(y)$–curves have the familiar S–shape, representing transversally directed collective flow of matter. This matter corresponds to the deflected spectators of Fig. 5(b). In all three cases the maximum of the transverse momentum is slightly below beam rapidity (because the deflection of the spectators decreases their longitudinal momentum) and of the order of 300 MeV. The maximum decreases slowly as a function of beam energy.

On the other hand, for the EoS with phase transition the transverse momentum has an S–shaped form for low beam energies, Fig. 6(a), although the maximum transverse momentum is by about a factor of two smaller than for the case of a pure hadronic EoS. This is due to the creation of a small amount of mixed phase matter in the very central region at these energies (cf. Fig. 2)[4]. However, around 5 AGeV (Fig. 6(b)) there is no longer any appreciable amount of collective transverse flow, the transverse momentum is essentially zero as function of rapidity. As explained in connection with Fig. 5(a), spectators are not deflected at all in this case, they pass the participant matter before the latter is expanding. This expansion is then more or less isotropic. For even higher beam energies, Fig. 6(c), the directed flow of matter gradually starts to increase again, since QGP with a higher p/ϵ is created (cf. Fig. 2(e)), which does again expand (and cool) more rapidly, see Fig. 4, and consequently deflects spectator matter.

In Fig. 7 we show the excitation function of the *directed* transverse (in–reaction–plane) fluid momentum per baryon,

$$\langle p_x/N \rangle^{dir} = \frac{1}{N} \int_{-y_{CM}}^{y_{CM}} \mathrm{d}y \; \langle p_x/N \rangle(y) \; \frac{\mathrm{d}N}{\mathrm{d}y} \; \mathrm{sgn}(y) \; , \tag{6}$$

which is in principle nothing else but the integral over the curves of Fig. 6, weighted with the baryon number rapidity distribution (and the sign of the rapidity, otherwise momentum conservation would yield a trivial value). First of all, one observes that above $E_{Lab}^{kin} \simeq 2$ AGeV this quantity decreases for increasing beam energy for calculations with the pure hadronic EoS (dotted line). This is in accord with Figs. 6(a–c) and simply due to the fact that faster spectators are less easily deflected by the hot, expanding participant matter. Second, the directed transverse momentum as calculated for the EoS with phase transition (full line) shows a dramatic *drop* between BEVALAC and AGS beam energies as compared to the calculation with the pure hadronic EoS and *increases* again beyond ~ 10 AGeV. Thus, *there is a local minimum in the excitation function of the directed transverse (in–reaction–plane) collective flow around* ~ 6 AGeV, which is related to the phase transition to the QGP and the existence of a "softest point" in the nuclear matter EoS. We note that the position of the minimum strongly depends on the EoS (cf. discussion of Fig. 2). It may easily shift to higher beam energies, if resonances are included in the hadronic part of the EoS. Also, as was the case for Fig. 6, absolute values for the directed momentum cannot be compared to experimentally measured ones, since a freeze-out calculation has not yet been performed. Such a quantitative comparison is in any case not reasonable at this stage since viscosity effects are neglected in the ideal hydrodynamic picture, which are known to have a strong influence on flow [35][5].

[4] We note that for central cells the compression is near the limiting values given by the one-dimensional collision scenario as shown in Fig. 2.

[5] How to implement dissipative effects into relativistic hydrodynamics is a presently unsolved problem,

The main point is, however, that irrespective of these quantitative uncertainties, the *minimum* is a clean *qualitative* signal for a transition from hadron to quark and gluon degrees of freedom in the nuclear matter EoS. We even expect this signal to be independent on whether the transition is a first order phase transition or merely a rapid increase of the energy density as a function of temperature (or, for baryon-rich matter, of the baryo-chemical potential). Let us emphasize that the mere fact that the flow *decreases* when using an EoS with phase transition relative to the pure hadronic case, as was found in [9], is not sufficient to serve as an unambiguous QGP signature. First of all, absolute values for the flow are not reliable due to the above mentioned uncertainties in the ideal hydrodynamic picture. Second, as can be seen in Fig. 7 such a decrease occurs also for the pure hadronic scenario and is due to trivial kinematic reasons. Finally, it was suggested in [14] that an observed flow which is less than predicted by cascade models might be a signature for QGP formation. This, however, appears to be insufficiently unique as well: as was shown in [13, 36] there is considerable freedom in treating two–particle scatterings in a cascade which results in quite different values for the flow.

In order to observe the minimum in $\langle p_x/N\rangle^{dir}(E_{Lab}^{kin})$ experimentally, it is mandatory to supplement present data on collective flow for BEVALAC [6] and AGS energies [37] with data taken at beam energies between BEVALAC and present AGS energies, in order to map out the excitation function of the transverse collective flow. Such experiments are currently under way at the AGS [38] and will help to decide whether physics at these energies is satisfactorily described by hadronic interactions, like in microscopic models [12, 13, 14], or whether one has already entered the rather interesting domain of qualitatively new phenomena that can only be attributed to the presence of quark and gluon degrees of freedom. We conclude by remarking that there are no data that conclusively exclude the latter possibility. Rather, the measured baryon rapidity distribution at AGS energies [39] can be equally well described by microscopic [13] and fluid-dynamical models [9].

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Figure Captions:

Fig. 1: The nuclear matter EoS in the form $p(\epsilon, n)$. Pressure and energy density are normalized to the ground state energy density $\epsilon_0 \simeq 146.5\,\mathrm{MeV\,fm^{-3}}$, the baryon number density to the ground state density $n_0 \simeq 0.16\,\mathrm{fm^{-3}}$.

Fig. 2: Excitation functions of various thermodynamic quantities for the one-dimensional "slab-on-slab" collision scenario. (a) ϵ/ϵ_0, (b) n/n_0, (c) p/ϵ_0, (d) σ, (e) p/ϵ, (f) T, (g) the velocity of sound squared c_s^2, and (h) the volume fraction of QGP λ in the final compressed state. Full lines are for the EoS with phase transition, dotted lines for the pure hadronic EoS. Regions of pure hadronic, mixed, and pure quark matter in the final compressed state are indicated by thin dotted lines.

Fig. 3: The time t_F a compression front requires to traverse an incoming nucleus in units of the nuclear rest frame radius R as a function of the beam energy. The full line is for the EoS with phase transition, the dotted for the pure hadronic EoS.

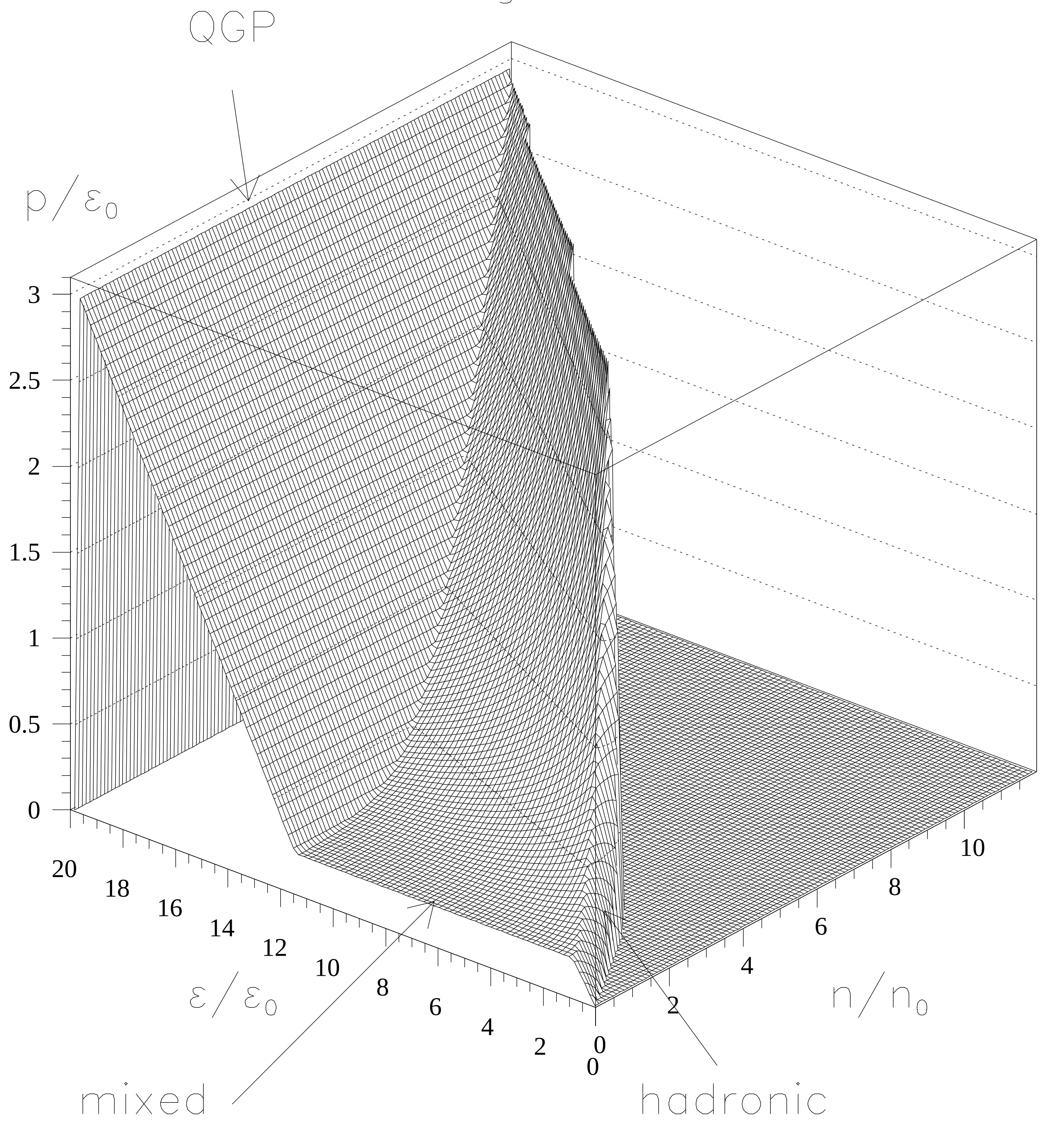
Fig. 4: The lifetime of the mixed phase, defined as the CM time when the T_c–isotherm intersects the t–axis at $x = 0$, in units of the radius of the hot, compressed system as a function of the initial energy density (in units of the phase transition pressure p_c). ϵ_H and ϵ_Q indicate the phase boundaries between hadron matter and mixed phase as well as mixed phase and QGP. The full line is for the EoS with phase transition, the dotted line for the pure hadronic EoS, assuming $T_c \simeq 170$ MeV.

Fig. 5: CM frame baryon density contours and flow velocity vectors in the reaction plane ($x - z$–plane) at different CM times for a 5 AGeV Au+Au–collision at $b = 3$ fm calculated with (a) the EoS with phase transition and (b) the pure hadronic EoS.

Fig. 6: The mean transverse (in–reaction–plane) momentum per baryon as a function of longitudinal (fluid) rapidity for Au+Au–collisions at $b = 3$ fm and for (a) $E_{Lab}^{kin} = 3.5$ AGeV, (b) 5 AGeV, and (c) 11.7 AGeV. The crosses correspond to calculations using the EoS with phase transition, the open circles to those with the pure hadronic EoS. Full and dash–dotted lines are interpolations between calculated points to guide the eye.

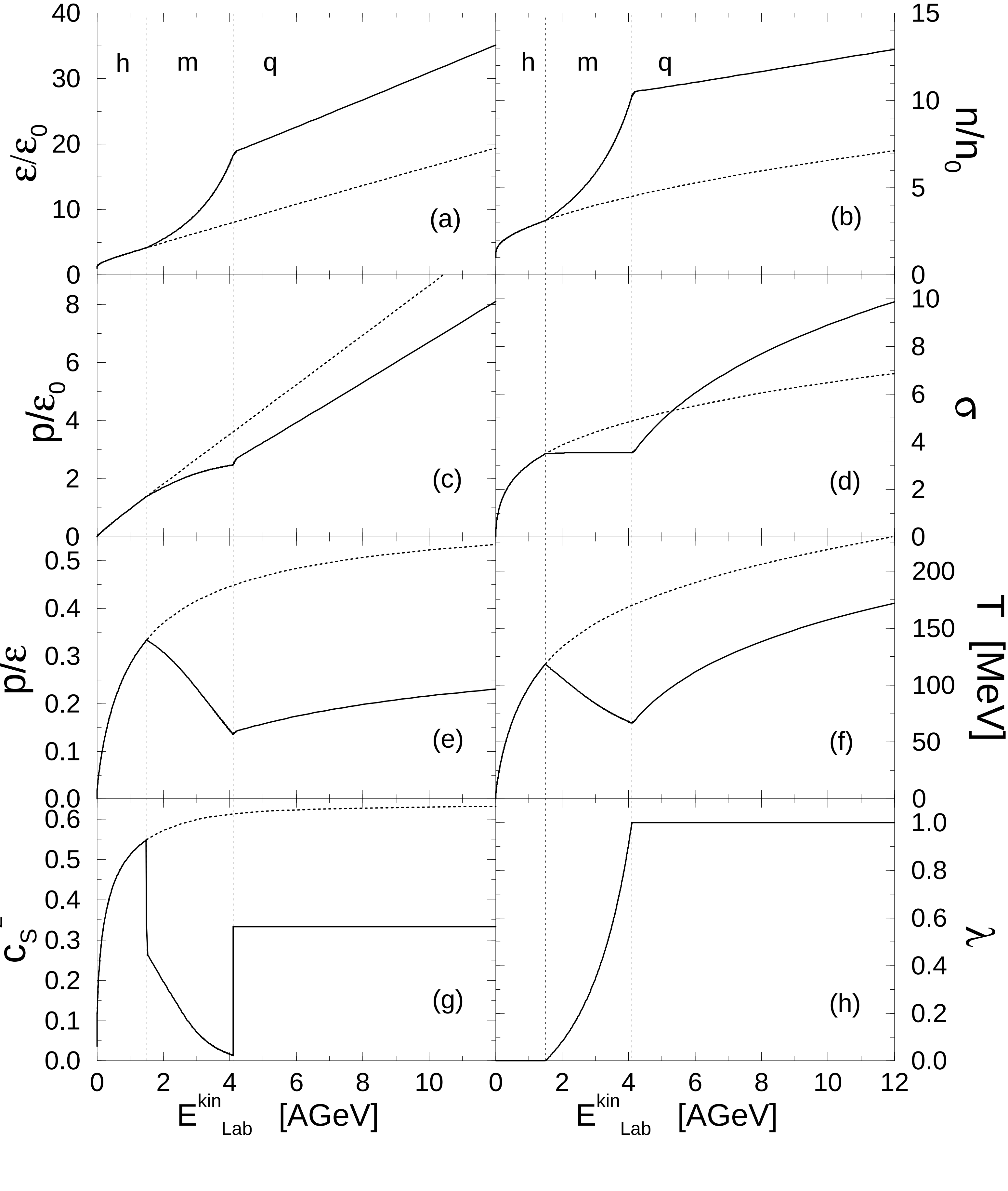
Fig. 7: The directed mean transverse (in–reaction–plane) momentum as a function of beam energy for Au+Au–collisions at $b = 3$ fm. The full line (crosses) corresponds to calculations using the EoS with phase transition, the dotted line (open circles) to those with the pure hadronic EoS.

Fig. 1



QGP
p / ε0
3
2.5
2
1.5
1
0.5
0
20
18
16
14
12
10
8
6
4
2
0
ε / ε0
0
2
4
6
8
10
n / n0
mixed
hadronic

Fig. 2



h
m
q
(a)
(b)
(c)
(d)
(e)
(f)
(g)
(h)
ε/ε0
n/n0
p/ε0
σ
p/ε
T [MeV]
cs 2
λ
E kin Lab [AGeV]
E kin Lab [AGeV]

Fig. 3

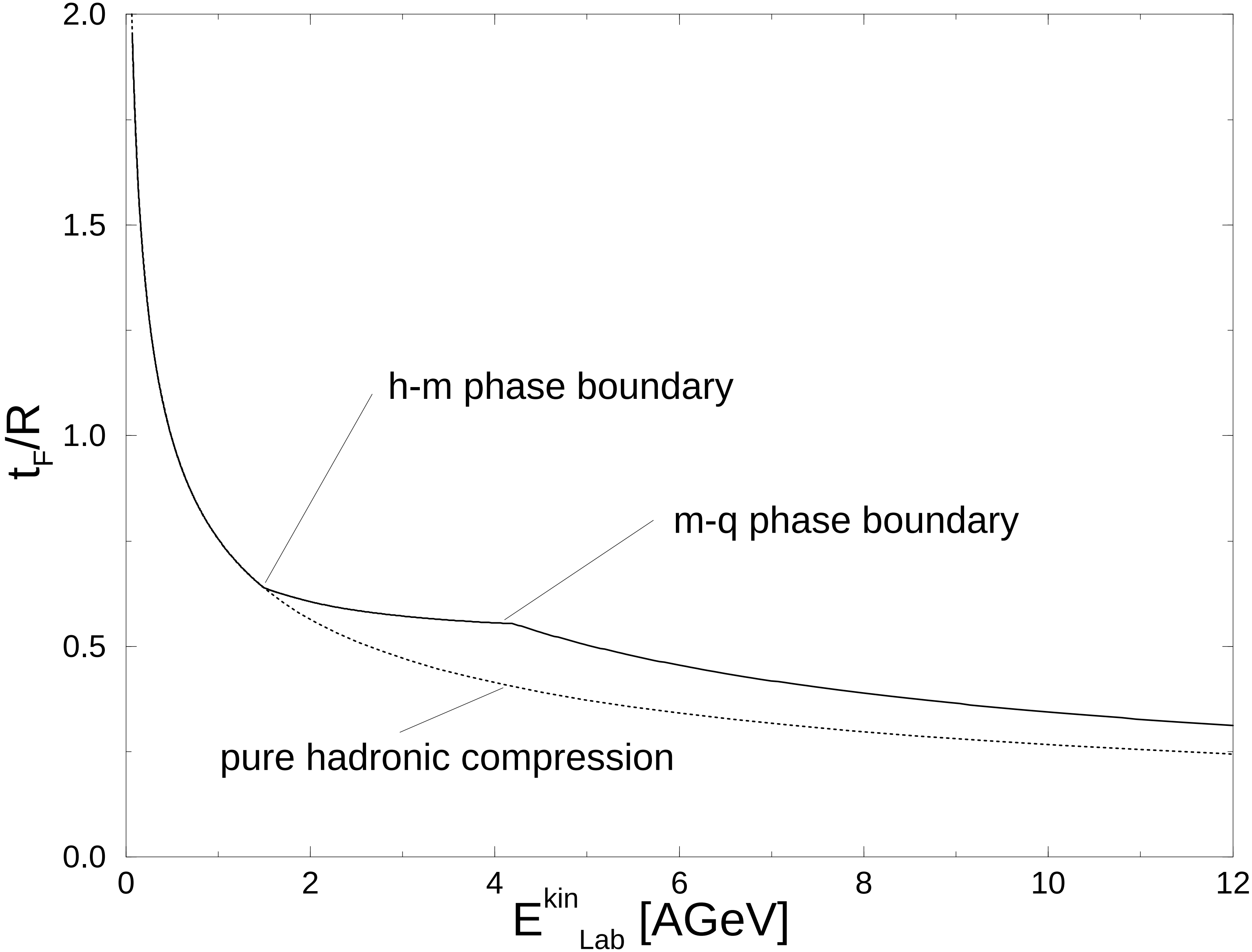
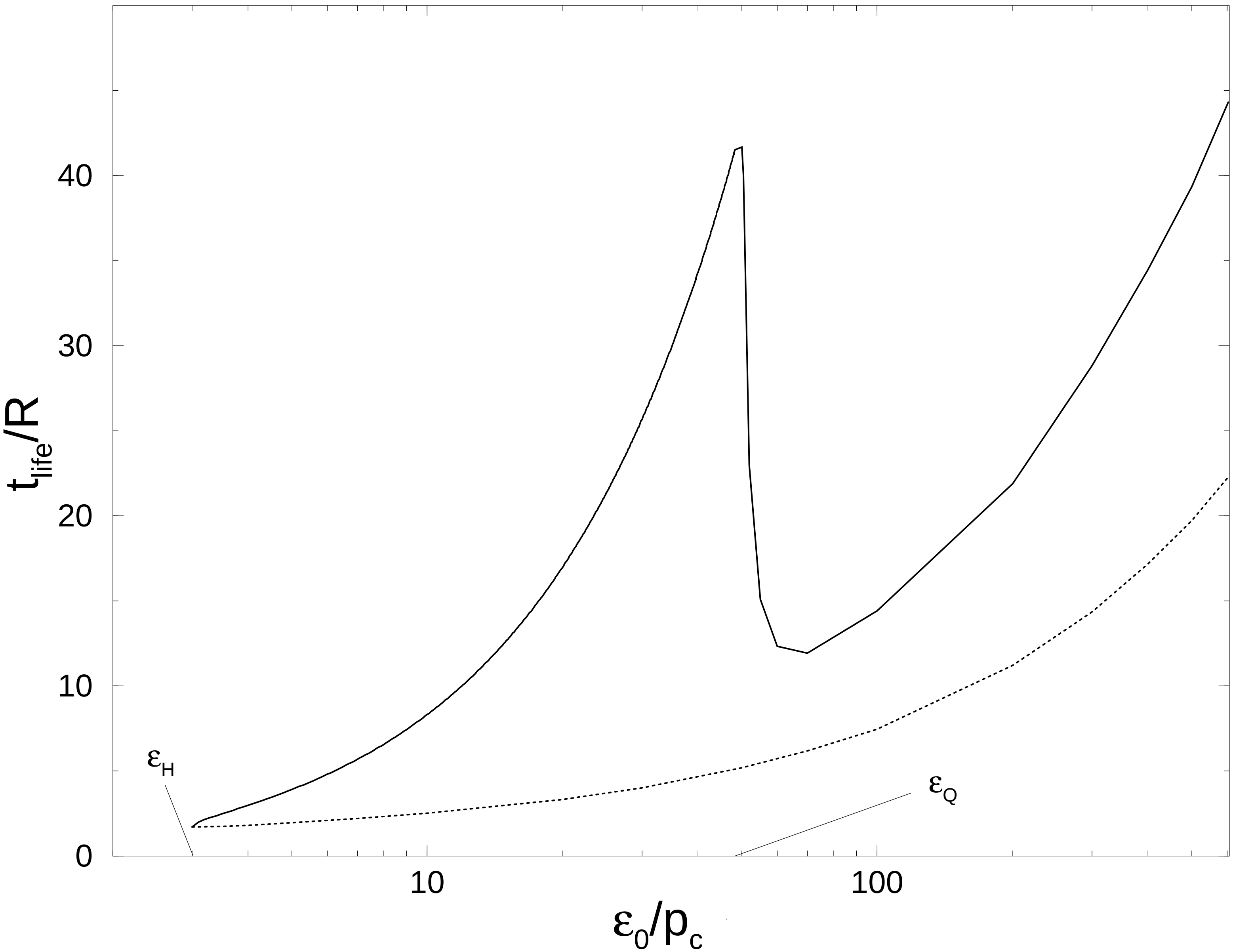

2.0
1.5
1.0
0.5
0.0
t_F/R
h-m phase boundary
m-q phase boundary
pure hadronic compression
0
2
4
6
8
10
12
E^{kin}_{Lab} [AGeV]

Fig. 4



t_{life}/R
40
30
20
10
0
ε_H
ε_Q
10
100
ε_0/p_c

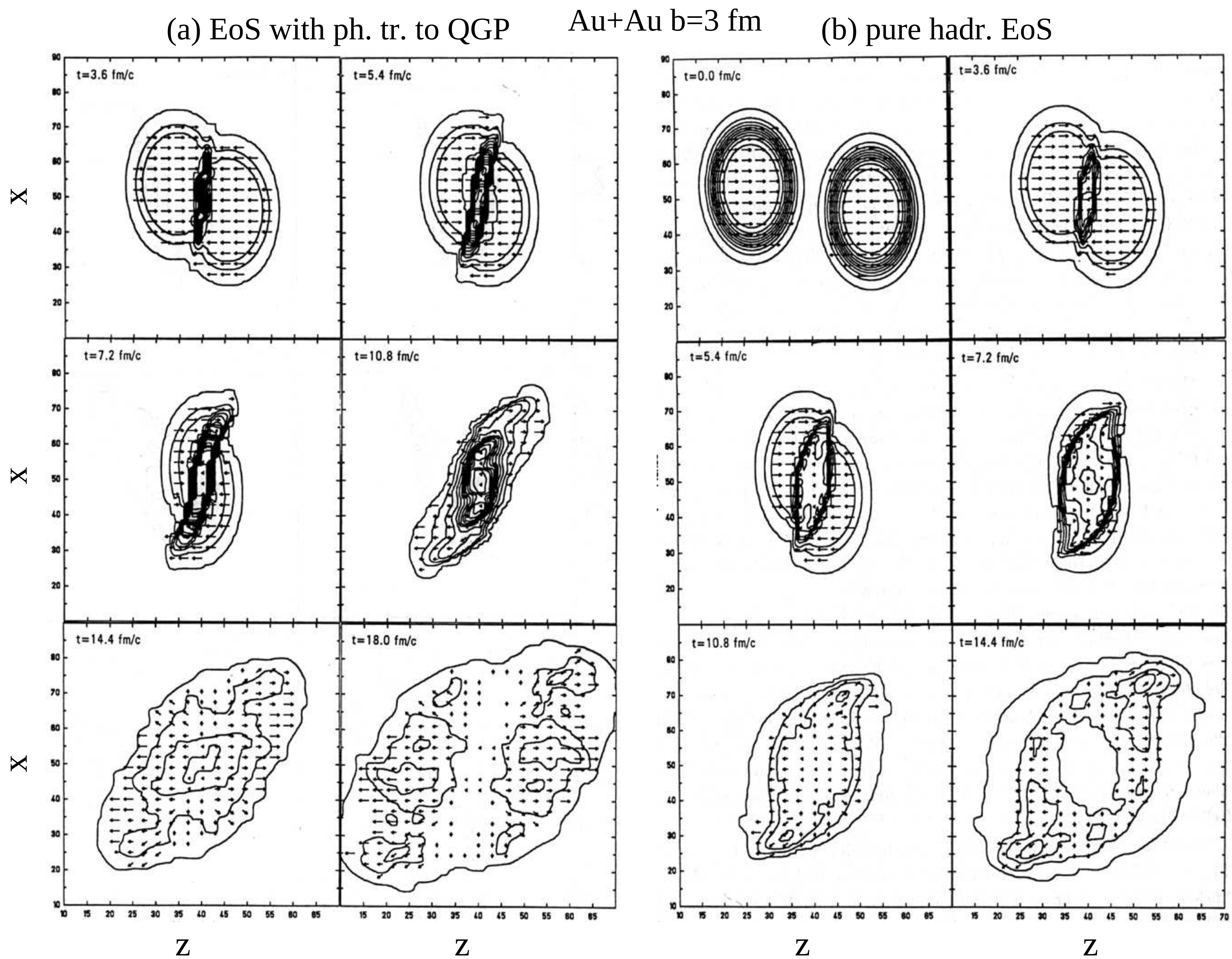

(a) EoS with ph. tr. to QGP
Au+Au b=3 fm
(b) pure hadr. EoS
t=3.6 fm/c
t=5.4 fm/c
t=0.0 fm/c
t=3.6 fm/c
t=7.2 fm/c
t=10.8 fm/c
t=5.4 fm/c
t=7.2 fm/c
t=14.4 fm/c
t=18.0 fm/c
t=10.8 fm/c
t=14.4 fm/c
X
Z

Fig. 5

Fig. 6

Au+Au b=3 fm

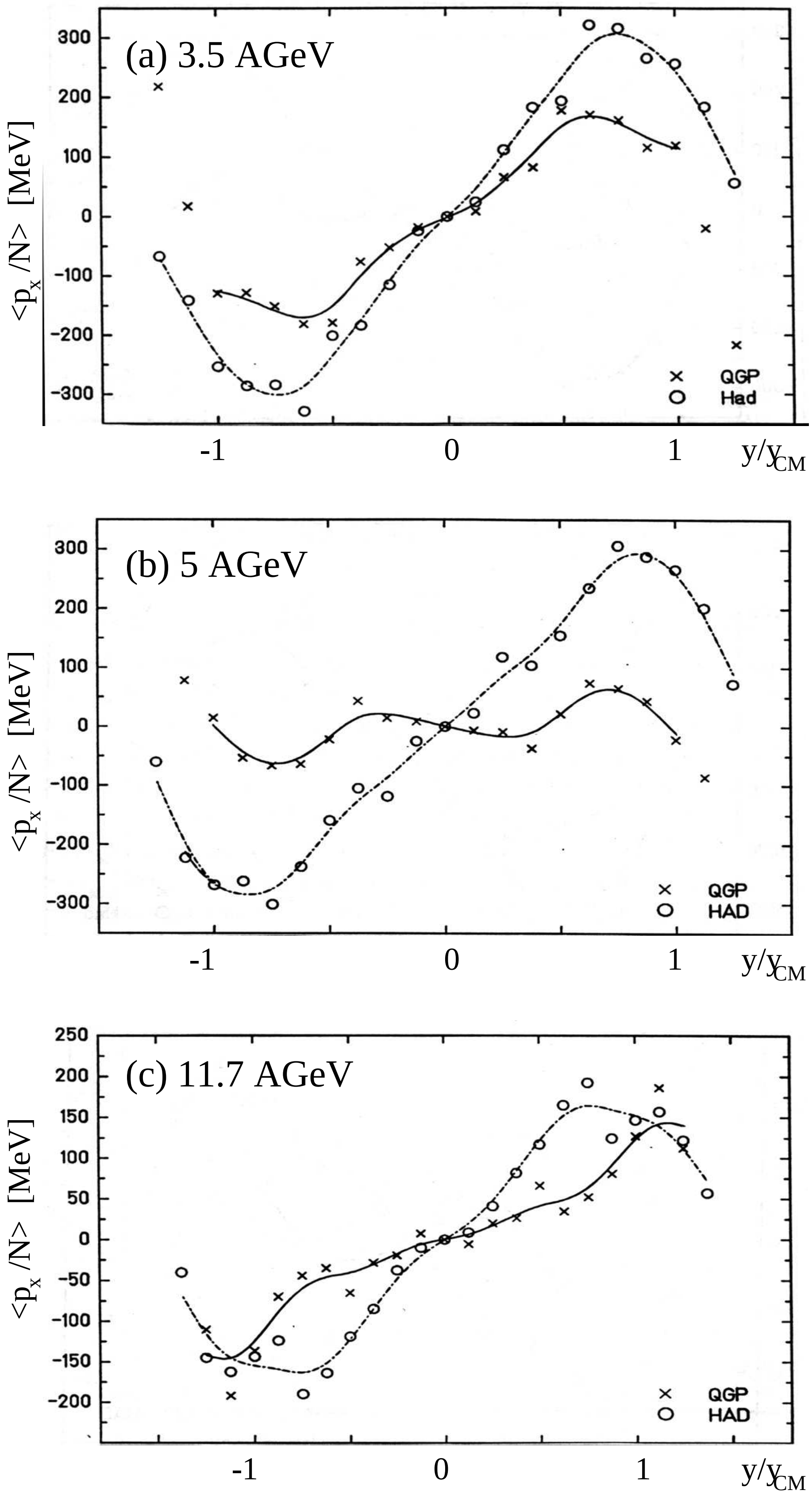
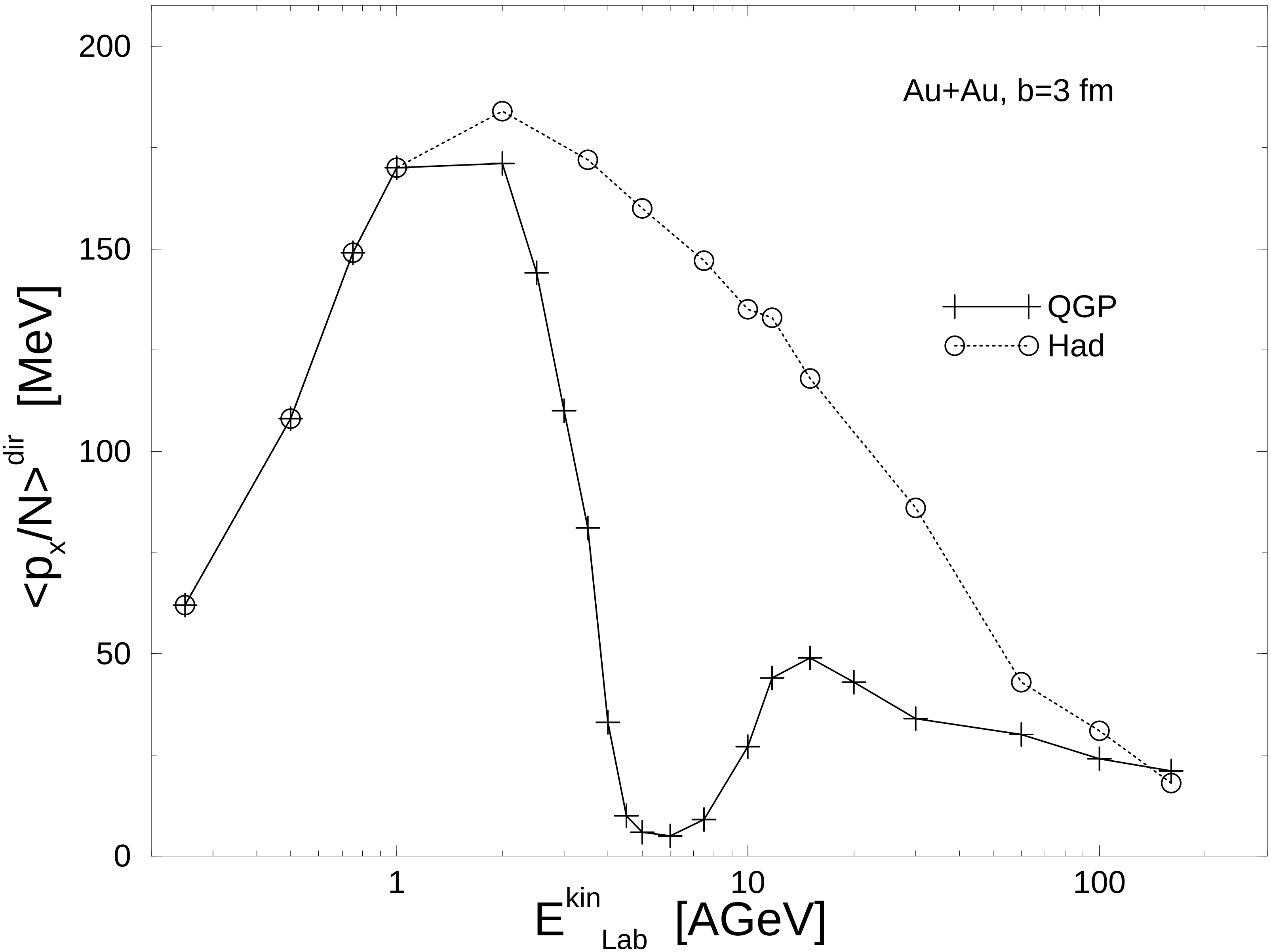


Fig. 7



Au+Au, b=3 fm
QGP
Had
200
150
100
50
0
1
10
100
<p_x/N>^dir [MeV]
E^kin_Lab [AGeV]