DECONFINEMENT OF CONSTITUENT QUARKS AND THE HAGEDORN TEMPERATURE

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INTRODUCTION

The overwhelming majority of theoretical papers on phase transition of hadronic matter $H$ to quark gluon plasma ($H \leftrightarrow QGP$) ignores constituent quarks $Q$ which we shall call below briefly valons (the term proposed by R.Hwa). The present paper is fundamentally based on the conception of those valons as real entities having the same quantum numbers as current quarks and the mass which was more than once calculated theoretically as some 300 MeV and used for explanation of experiments (e.g., see\(^1\)). The possibility of existence of a special state of strong interacting matter, that of deconfined valon gas with still broken chiral symmetry, is discussed here.

Such a possibility, i.e. appearance of the third, $Q$ phase intermediate between $H$ and $QGP$ had been discussed already in papers\(^2\)–\(^5\) (we apologize for possibly missing some other papers). The idea may be traced back to E.Shuryak who pointed out that deconfinement and restoration of chiral symmetry might not coincide\(^6\). However in the papers mentioned above the result turned out pessimistic. E.g., on the $\mu - T$ diagram ($\mu$ - chemical potential, $T$ - temperature) possible $Q$ phase (allowing also for admixture of pions necessary first of all as Goldstone particles) occupies but a very small area at large $T$ and rather small $\mu$ (see e.g. Fig.1 taken from\(^5\)). The ”constituent quark phase is strongly suppressed”\(^2\).

Meanwhile the intuitive physical pattern\(^7,8\) based on the bag model ideology, although being extremely rough, points to the quite a different possibility schematically shown in Fig.2. In fact, a nucleus consists of nucleons $N$ which in turn are bags containing valons, which are nothing else than a massive clouds of strong interacting current quarks $q$, antiquarks $\bar{q}$ and gluons $g$ with quantum numbers of a single quark. If we compress the nucleus (Fig.2a) until nucleons become close packed and fill up the entire space (Fig.2b) they merge into a single bag with valons deconfined (still with chiral...
symmetry broken) and freely moving within the common bag (Fig. 2c). Further compression leads to a new situation, that of close packed and merging valons (Fig. 2d). Now \( q, \bar{q} \) and \( g \) which earlier were confined within valons can freely propagate within a common bag (Fig. 2e). Here they are deconfined and massless, chiral symmetry is automatically restored and this very state is \( \text{QGP} \).

Thus we expect two phase transitions to take place: the one, \( H \leftrightarrow Q \), and (at higher density and/or temperature) another, \( Q \leftrightarrow \text{QGP} \).

In this paper we are going to show, in distinction to \( 2^{-5} \), that with reasonable values of bag parameters \( B_Q \) for the valon phase and \( B_q \) for \( \text{QGP} \) there actually appears possibility for a pattern shown in Fig. 2, i.e. the possibility of the overall \( H \leftrightarrow \text{QGP} \) transition to proceed only via the intermediate \( Q \) phase.

The first \( (H \leftrightarrow Q) \) transition establishes limitations on possibility of existence of normal hadrons. This is the reason for interpreting it as the Hagedorn limit and to consider the temperature of transition, i.e. of valon deconfinement, \( T_d \), as the Hagedorn temperature \( T_H \). The second transition restores chiral symmetry and is characterized by the temperature \( T_{ch} \). Thus the Hagedorn temperature appears to be not merely a very rough approximation to chiral symmetry restoration/breaking temperature as is frequently believed. It has its own physical meaning of temperature of constituent quark (≡valon) deconfinement: \( T_H \equiv T_d \).

**THERMODYNAMICAL FORMALISM**

We are treating phase transitions following traditional manner \(^9_{2^{-5}}\), fundamentals of which were given by Hagedorn \(^9\). As the first step we consider the partition function \( Z^0(T, \mu, V) \) depending on \( T, \mu, \) and system volume \( V \) for each phase in the first approximation (which means: for pointlike particles):

\[
\ln Z_j^0(T, \mu, V) = \frac{V}{T} \sum_i \left\{ \frac{g_i^B}{6\pi^2} \int \frac{k^4 dk}{\sqrt{k^2 + m_i^2}} \exp \left( \frac{1}{T} \right) - 1 \right\} + \\
\frac{g_i^F}{6\pi^2} \int \frac{k^4 dk}{\sqrt{k^2 + m_i^2}} \left[ \frac{1}{\exp \left( \sqrt{k^2 + m_i^2} - \mu \right)} + 1 \right] + \ln Z_{vac}^0
\]

where \( g_i^B,F \) and \( m_i \) are degeneracy factors and masses for Bose and Fermi \( i \)-th type particles respectively for each \( j \)-th phase (\( j \) designates hadronic \( H \), valonic \( Q \), and current \( \text{QGP}, q \) phases respectively). \( \ln Z_j^0(T, \mu, V) \) is a result of summation over particles which play considerable role (under \( T \) and \( \mu \) in question) in the given phase. In particular, for hadronic phase \( H \) we have taken into account nucleons (protons and neutrons, \( M = 940 \text{ MeV} \)), pions (\( m_\pi = 140 \text{ MeV} \)) and \( \Lambda \) hyperons (\( m_\Lambda \approx 1150 \text{ MeV} \)). The constituent phase \( Q \) contains light valons \( \equiv \) constituent quarks (\( m_u \approx m_d \approx 320 \text{ MeV} \)), some admixture of strange valon quarks (\( m_s \approx 512 \text{ MeV} \)), and again pions. The term \( \ln Z_{vac}^0 \) stands here for \( Q \) and \( \text{QGP} \) phases and reflects effective interaction of quarks and gluons with the QCD-vacuum. This interaction in \( Q \) and \( \text{QGP} \) phases is assumed as

\[
\ln Z_{vac,Q}^0(T, \mu, V) = -\frac{V}{T} B_Q ,
\]

\[
\ln Z_{vac,q}^0(T, \mu, V) = -\frac{V}{T} B_q
\]

(1)
where \( B_Q, B_q \) are some constant bag parameters differing from each other (to be discussed below).

Then by usual differentiations of \( Z_0^j(T, \mu, V) \) we obtain for each phase energy density \( \epsilon_j^{(0)} \), pressure \( p_j^{(0)} \), and particle density \( n_j^{(0)} \) as functions of internal parameters \( T, \mu, V \).

Having this as a basis and following the procedure proposed \(^2\)–\(^5\) (more exactly in \(^2\)) we introduce mutual interaction of particles by ascribing them hard core radii in the Van der Waals manner. Namely, we substitute \( n_j^{(0)} \) with real particle density \( n_j^{(r)} \) and then obtain ”real” energy density \( \epsilon_j^{(r)} \) and pressure \( p_j^{(r)} \) for each of three phases \( j = H, Q, QGP \):

\[
\begin{align*}
n_j^{(r)}(T, \mu) &= \frac{n_j^{(0)}(T, \mu)}{1 + n_j^{(0)}(T, \mu) * v_j}, \\
\epsilon_j^{(r)}(T, \mu) &= \frac{\epsilon_j^{(0)}(T, \mu)}{1 + n_j^{(0)}(T, \mu) * v_j}, \\
p_j^{(r)}(T, \mu) &= \frac{p_j^{(0)}(T, \mu)}{1 + n_j^{(0)}(T, \mu) * v_j},
\end{align*}
\]

where each of the parameters \( v_j \) is hard core volume characterizing \( j \)-th phase particles:

\[
\begin{align*}
v_N &= \frac{4}{3} \pi r_N^3, \\
v_Q &= \frac{4}{3} \pi r_Q^3,
\end{align*}
\]

current quarks and gluons are of course pointlike \( (v_q = 0) \).

At given \( T \) and \( V \) of the system the most preferable (i.e. stable) phase is the one having the largest pressure \( p_j \) and smallest \( \mu_j \). At equilibrium of each two phases there pressures and effective chemical potentials are to be equal. Three types of phase equilibrium curves are to be analyzed: the one corresponding to the valon deconfinement transition \( (H \leftrightarrow Q) \)

\[
\mu_H = 3\mu_Q; \quad p_H(T_d, \mu_H) = p_Q(T_d, \mu_H/3),
\]

**direct** transition \( (H \leftrightarrow QGP) \)

\[
\mu_H = 3\mu_q; \quad p_H(T_c, \mu_H) = p_q(T_c, \mu_H/3),
\]

and **chiral** transition \( (Q \leftrightarrow QGP) \)

\[
\mu_Q = \mu_q; \quad p_Q(T_{ch}, \mu_Q) = p_q(T_{ch}, \mu_H/3).
\]

The point of **coexistence of all three phases** (i.e. the **triple point**) is defined by the following conditions:

\[
p_H(T^#, \mu^#) = p_Q(T^#, \mu^#) = p_q(T^#, \mu^#)
\]

Before turning to study of equilibrium conditions and phase transitions let us discuss the chosen set of parameters.

**BASIC PARAMETERS**
There are two pairs of decisively important parameters: firstly, particle sizes, e.g. average root square radii of nucleons $N, r_N$, understood as $(\bar{r}_N^2)^{1/2}$, and valons, $r_Q$, which designates $(\bar{r}_Q^2)^{1/2}$. Secondly bag pressures, $B_Q$ for the $Q$ phase and $B_q$ for QGP.

Let us discuss sizes. They are used to describe interaction of particles in a simplified form by considering them as black balls forming a Van der Waals gas. This means that we approximate their interaction potential $U(r)$ as $U = \infty$ for $r < r_c$ and $U = 0$ for $r > r_c, r_c$ being the core radius (in other calculations we have used also some continuous potential $U(r)$). In the nucleon interaction was instead described by two versions of the continuous $U(r)$ satisfactorily describing nuclear properties. It seems sufficient in this paper to consider $N$ and $Q$ simply as having cores. For nucleon radius and volume $r_N, v_N$ we usually assume, as most reasonable, the values:

$$r_N \simeq 0.8\text{fm}; \quad v_N \simeq 2.13\text{fm}^3$$

(however sometimes we shall consider also $r_N = 0.7$ fm). Since the nucleus has the radius $R_A = r_0 \ast A^{1/3}$ ($A$ - atomic weight) with $r_0 \simeq 1.10 - 1.15$ fm, the fraction of the nucleus volume occupied by nucleons is

$$(r_N/r_A)^3 \simeq 0.26 \div 0.38.$$  \hspace{1cm} (14)

Thus it seems sufficient to compress a nucleus only 3-4 times to come to close packing of nucleons.

The size of valons has been estimated from analysis of properties of various hadrons within the Additive Quark Model (AQM); when expressed by the square root radius $r_Q$ (for $r_N = 0.8$ fm) this estimate sounds as:

$$0.20 \leq r_Q \leq 0.36\text{fm}.$$ \hspace{1cm} (15)

We shall concentrate on some middle reasonable value, $r_Q = 0.3$ fm, although other values within limitation (15) will be also used (let us remark that estimating $r_Q$ within black ball valon model and AQM from $QQ$ collision cross section $\sigma_{QQ} = \pi(2r_Q)^2 \simeq (1/9)\sigma_{NN}$ where $\sigma_{NN}$ is the $NN$ collision cross section taken, at not too high energies, as 40 mb, we get $r_Q = 0.2$ fm. For tevatron energies it is much larger).

Now, about $B_Q$ and $B_q$. The choice of their values is much more disputable, being at the same time very essential. Previous authors were (seemingly) guided by the results of lattice calculations in which confinement/deconfinement $(c/d)$ and chiral symmetry breaking/restoration $(b/r)$ transitions had been found to coincide at $\mu = 0$. Moreover ”deconfinement” was (and still is) usually understood as the straightforward $H \leftrightarrow QGP$ transition. Accordingly, e.g. a condition was superimposed on pressures $p_H$ and $p_q$ of $H$ and QGP phases at the transition point (for $\mu = 0$) : $p_H(T_c, 0) = p_q(T_c, 0)$. Herefrom the ratio $\beta = B_q/B_Q$ was obtained thus leaving only one of these two parameters undefined. In other works similar coincidence of $(c/d)$ and $(b/r)$ at $\mu = 0$ was also taken as a basis. Herefrom, for a freely chosen $B_Q$, followed the value of $B_q$. All authors accordingly used mainly the values $\beta = B_q/B_Q \sim 3/4$ with $B_Q$ more or less close to the MIT bag value, $B_{MIT} \sim 56$ MeV/fm$^3$ ($B_Q = 56 \div 150$ MeV/fm$^3$, $B_q = 200 \div 500$ MeV/fm$^3$, keeping in all cases $\beta \sim 3/4$).

Assumption concerning $B_Q \simeq B_{MIT}$ seems to be reasonable at least in the vicinity of the $H \leftrightarrow Q$ transition. In fact, here valon density in $Q$ phase, at least in the beginning of the valon deconfinement process, is the same as in a nucleon (see Fig.2b). Accordingly we shall use the assumption:

$$50 \leq B_Q \simeq B_{MIT} \leq 100 \text{ MeV/fm}^3.$$ \hspace{1cm} (16)
On the contrary, argumentation which have lead to the above mentioned estimate of $\beta = B_q/B_Q \sim 3$ is not admissible for us since difference between confinement/deconfinement (c/d) of valons (which takes place with broken chiral symmetry) and chiral symmetry breaking/restoration (b/r) is exactly the phenomenon we are looking for. Accordingly this difference should not be excluded at the start of the analysis as has been, as a matter of fact, done e.g. in\textsuperscript{5}. Thus $B_q$ is a parameter which needs its own physical foundation. Since it refers to QGP, it should be tightly connected with the vacuum pressure $B_{vac} \simeq 500 \div 1000$ MeV/fm$^3$ and be close to it. Accordingly we are going to use larger values of $\beta = B_q/B_Q$ putting it equal to $\beta \sim 10$. Below we shall use the following set of reasonable parameters as the standard one:

$$B_q = 500 \text{ MeV/fm}^3; \; B_Q = 50 \text{ MeV/fm}^3; \; r_N = 0.8 \text{fm}; \; r_Q = 0.3 \text{fm}.$$  \hspace{1cm} (17)

The influence of $B_Q$ and $B_q$ as well as $r_j$ variations are partially investigated.

**MAIN RESULTS**

Let us start with extreme cases: $\mu = 0, T \neq 0$ (Figs.3,4) and $T = 0, \mu \neq 0$ (Fig.5).

Fig.3 clearly demonstrates characteristic pattern of double phase transition at the $p - T$ plane for chosen standard parameter set (17). It is evident that two phase transitions are well separated, $Q$ is the most stable one (i.e. its pressure $p_Q$ is higher than that of other phases) within the temperature interval of $\Delta T = T_{ch} - T_d \sim 50$ MeV, with absolute values of transition temperatures being of the order of the direct phase transition ($H \leftrightarrow \text{QGP}$) temperature $T_c$: $T_c \simeq 180$ MeV, $T_d \simeq 145$ MeV, $T_{ch} \simeq 195$ MeV. The value of this "temperature corridor" of Q phase depends on $B_q$ (see Fig.4(b)) so that it becomes larger (up to 90 MeV) for larger values of $B_q$ (when $B_q = 1$ GeV/fm$^3$, $\beta \sim 20$, $T_{ch} \sim 230$ MeV) and shrinks to a triple point (which are marked by the symbol "#" within all the figures) $T^\# \sim 140$ MeV for most frequently choosen $B_q \simeq 200$ MeV/fm$^3$ $\beta \sim 4$, close to that in\textsuperscript{5}. The same effect stays for $B_Q$ dependence: the closer it is to $B_q$, the narrower is $T$-region allowing for $Q$ phase existence, with the biggest value (corresponding to the triple point $T^\# \sim 180$ MeV) $B_Q \sim 160$ MeV (i.e. $\beta \sim 4.5$). For all combinations of $B$'s the chiral symmetry restores at $T_{ch}$ which is always larger than the critical temperature of direct transition $T_c \equiv T_{H \leftrightarrow \text{QGP}}$ obtained in calculations when ignoring existence of valons. Note that we do not show the influence of $r$ parameters because it is negligeable in this $T - \mu$ region.

The similar pattern is presented in Fig.5 for $T = 0, \mu \neq 0$, for the standard parameter set(17). Here again there is a possibility for appearance of $Q$ phase region but now it is greatly suppressed: if $\beta = 10$ while $B_Q$ has its minimal reasonable value 50 MeV/fm$^3$, the area covered by $Q$ phase on $\{p, \mu\}$ diagramm is rather small and it shrinks into a triple point already at $B_Q \simeq 80$ MeV/fm$^3$; this region becomes of course broader for larger $\beta$ values. However it is essential that valons tell even at $T = 0$. For more details concerning $T = 0$ see \textsuperscript{11}.

Now we go over to the general case of $T, \mu \neq 0$. The $\{\mu, T\}$ diagramm for the standard parameter set (17) is shown in Fig.6. It is clear that there exists a wide corridor around the curve for direct $H \leftrightarrow \text{QGP}$ transition which remains as well for various combinations of parameters ($B_q, B_Q, r_N$, and $r_Q$) (see Figs.7). It shrinks and even disappears with $\mu$ growth only at rather definite conditions, namely for small $\beta$ (the shape of $Q$ phase region for $\beta < 3$ shown in Fig.7(a), as well as in Fig.1, reminds a thin banana) and for extremely large $r_Q$ (Fig.7(b)). Otherwise transition from $H$ phase to QGP proceeds only via valon phase (for details see \textsuperscript{12}).
Note that under not too large $\mu \leq 1$ GeV the master parameter is $\beta = B_q/B_Q$ while influence of particle interaction (i.e. choice of values of $r_N$ and $r_Q$) is quite negligible. Symbol line in Fig.7(b) corresponds to the case of all particles taken point like. It differs considerably from the real particle diagram only at $\mu \sim 1$ GeV, and covers the area having the shape of some exotic fruit ("pineapple").

Critical values $\epsilon^\mu$ for energy densities (that is the ones on phase equilibrium curves) are presented in Figs.8. Let us point out (see Fig.8(a)) that the latent heat, i.e. the jump $\Delta \epsilon^d = (\epsilon^d_Q - \epsilon^d_H)$ at the deconfinement transition is much smaller than that at the chiral one, $\Delta \epsilon_{ch} = (\epsilon_{ch}^d - \epsilon_{ch}^Q)$, and much smaller than that for direct transition ignoring valons (Fig.8(b)). The deconfinement transition requires rather "soft" experimental conditions.

**DISCUSSION AND SUMMARY**

Of course, obtained numerical values, in particular those of $T_d$ and $T_{ch}$, should not be taken literary. E.g., they essentially depend on $B_Q$ and $B_q$ values chosen in a rather rough manner. According to e.g. Shuryak $^{13}$, for hadronic substance at $T = 0$, $B_{eff}$ (efficient $B_{eff}$ describing the summarized influence of interaction of particles between themselves and the vacuum) is not reducible to two constants but changes continuously increasing with baryon number density $n_B$ from some lower value of the order of $B_{MIT}$ up to $B_{eff} \sim B_{vac}$ for $QGP$. However, our calculations $^{12}$ show that baryon number density in $Q$ phase, $n_Q$, is much closer to that in the $H$ phase than to the one in the $QGP$. Thus the use of only two values - $B_Q$ and $B_q$ - seems to be not too bad.

Further, our calculations show that particle density of $H$ phase in the vicinity of transition curve to $Q$ phase at relatively large $\mu > 1$ GeV is very close to its limiting value $\sim v_N^{-1}$ (close packing) and use of the Van der Waals gas approximation for it is very suspicious. But the same weak point is met in the papers which lead to pessimistic conclusions concerning possibility of the $Q$-phase$^{2-4}$. Repulsive potential for nucleons considered in $^5$ seems of course to be more realistic; similar potential will be discussed in our paper $^{12}$. This does not change the final result qualitatively.

Nevertheless qualitative and even semiquantitative results seem to deserve attention. Comparing Figs.6,7 with Fig.1 we see that for sufficiently large values of $B_q/B_Q$ the pattern declared in the beginning of this paper (Fig.2) is almost fully supported. We may conclude that there actually exists possibility of hadronic phase to $QGP$ transfer (at seemingly reasonable choice of parameters) to proceed only via an intermediate valon, i.e. constituent quark plus pions, phase. Therefore, we expect existence of two phase transitions. For $\mu = 0$ transition temperatures differ by some 50 MeV. One of them, deconfinement of valons, proceeds at the Hagedorn temperature $T_d = T_H$, above which, in the deconfined valon phase, hadrons exist but as a small thermal admixture. The second phase transition at $T_{ch} \sim 200$ MeV corresponds to deconfiment/confinement of current quarks and gluons and, simultaneously, restoration/breaking of chiral symmetry. Thus difference between $T_d$ and $T_{ch}$ is physically meaningful.

The $Q$ phase with deconfined constituent quarks (valons) having mass $\sim 320$ MeV, due to $m_Q/T > 1$, is nonrelativistic, its state equation is that of nonrelativistic Van der Waals gas (since $r_Q$ is small). It is only in the vicinity of the phase transition to $QGP$ that valon size begins to tell. However, this last transition at not too large $\beta = B_q/B_Q$ takes place long before close packing of valons is attained$^{12}$. This result is rather interesting and has more physical sense than had been put into the model. In fact the model of hard black balls, being quite sufficient for nucleons due to their high
stability, is hardly valid for constituent quarks which may become partially transparent and larger. Overlapping of valon tails securing exchange of current quarks and gluons between valons may happen at distances exceeding the assumed black ball radius. Thus the transition to QGP phase does not necessarily require for close packing of valons. This is not so for nucleons at $H \leftrightarrow Q$ transition where the transition happens long after the considerable close packing is attained; this may be ascribed to necessity of overcoming the surface tension of nucleons. However for more detailed analysis it is desirable to consider more realistic repulsive potential for interaction of valons than simple hard core model (e.g. at the same manner as it has been done in\textsuperscript{14}).

The obvious first application of these results of course is to be to the evolution of matter formed at collision of highly relativistic heavy nuclei. Let us stress two most important items.

First, as it have been shown above (Figs.8) the latent heat value for $H \leftrightarrow Q$ transition is much smaller than that for the "upper" transition $Q \leftrightarrow QGP$ (and much smaller than people had obtained for direct, i.e. $H \leftrightarrow QGP$ transition). The same is true\textsuperscript{12} for transition pressure and $n_B$ values. This means that first deconfinment transition may proceed under much more "soft" conditions than it was expected for direct transition $H \leftrightarrow QGP$: it is sufficient to compress the nucleus only 3-4 times applying not very high pressure. Thus this transition may happen at not too high energies of colliding nuclei, i.e. even at Bevalac or in Dubna\textsuperscript{7,8} where appearance of pure QGP is hardly possible. Accordingly, e.g.in the Dubna experiments one may expect to observe some effects impossible for simple $pp$ collisions. In particular, our estimates show that admixture of $\Lambda$ particles may attain in central rapidity region some 10 percents of nucleons and this reminds us Okonov group Dubna experiment\textsuperscript{15} in which $\Lambda$'s were especially studied.

Secondly, transition pattern described above should influence essentially the character of hydrodynamical expansion and cooling of initially hot matter both if it takes place either in $Q$ or in $QGP$ phase. For $QGP$ initial state the following picture seems to be natural. The system experiences the phase transition to $QGP$ phase. After cooling down to chiral transition temperature $T_{ch} \sim 200$ MeV, it would further expand and cool down within a large temperature interval according to nonrelativistic state equation. Since at present ( and later when RHIC will be operating ) the attainable initial temperature may exceed $T_{ch}$ but slightly, the time spent for getting $T_{ch}$, as well as duration of cooling down in $H$ phase from $T_d \sim 150$ MeV to freezing temperature, $T_f \sim 140$ MeV, may be smaller than the time of existence of $Q$ phase with its nonrelativistic state equation. Therefore the existence of $Q$ phase should tell itself essentially. Here detailed calculations for resulting rapidity distribution, etc. are necessary to compare with experiment (at present they are in progress).

The results obtained seem to contradict to the statement based at lattice calculations that at $\mu = 0$ two phase transitions are to coincide. The situation here is however disputable. Although existence of $Q$ phase was seemingly never especialy looked for in lattice calculations there are some works in which existence of "heavy modes" was assumed in order to get rid of some results looking physically unsatisfactory\textsuperscript{16} ( the QGP phase does not have properties of an ideal relativistic gas ). Lattice specialists should answer the question whether the double phase transition with $Q$ mass at 320 MeV could be noticed in lattice calculations performed until now and how it should show itself.
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FIGURE CAPTION.

Figure 1. Phase transition curves in the $T - \mu$ plane (from$^5$): $Q$-phase region is between dashed and dotted lines.

Figure 2. Schematic picture of nuclear matter transformation under compression.

Figure 3. Phase diagram for $\mu = 0$ case at the $P - T$ plane for standard parameter set(17): the solid, dashed and dotted lines correspond to the the pressures of hadronic ($H$), valonic ($Q$) and current $q(QGP)$ phases respectively.

Figure 4. The same plot as in Fig.3 for standard parameter set (solid lines) and varying $B$’s: (a) $B_Q \simeq 100$ (dashed line), 160 MeV/fm$^3$ (dotted line); (b) $B_q \simeq 1000$ (dashed line), 200 MeV/fm$^3$ (dotted line);”#” marks the triple points.

Figure 5. Phase diagram in the $p - \mu$ plane for $T=0$ and standard parameter set; designations are the same as in Fig.3.

Figure 6. Phase transition curves in the $T - \mu$ plane: valon deconfinement (dashed line), direct transition ignoring valons (solid line), chiral transition (short dashed).

Figure 7. Transition curves bounding the $Q$-phase region: (a) - for standard $r_j$ and $B_Q$, $B_Q = 50$ MeV/fm$^3$ (solid lines), 100 MeV/fm$^3$ (long-dashed lines) and 150 MeV/fm$^3$ (dotted lines); short-dashed line corresponds to $B_q = 1$ GeV/fm$^3$; (b) - for standard $B$’s and varying $r_Q = 0.2$ fm (solid lines), 0.4 fm (dashed lines) and $r_j = 0$ ($\dagger$ lines).

Figure 8. Energy densities $\epsilon_T^j$ curves for: (a) $H$ and $Q$ matter at deconfinement transition (solid lines), $Q$ and $QGP$ matter at chiral transition (dashed lines); (b) $H$ and $QGP$ matter at direct transition ignoring valons (dotted lines).

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\( \varepsilon_j^{cr} \), GeV/fm\(^3\)
\( \epsilon_j^{cr}, \text{ GeV/fm}^3 \)

\( QCP \)

\( H \)