QCD AGAINST BLACK HOLES OF A STAR MASS ?

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Abstract

Along with compacting baryon (neutron) spacing in a neutron star (NS), two very important factors come into play side by side: the lack of the NS gravitational self-stabilization against shutting to black hole (BH) and the phase transition - color deconfinement and QCD-vacuum reconstruction - within the nuclear matter the NS is composed of. That is why both phenomena should be taken into account at once, as the gravitational collapse is considered. Since, under the above transition, the hadronic-phase (HPh) vacuum (filled up with gluon- and chiral $q\bar{q}$-condensates) turns into the "empty" (perturbation) subhadronic-phase (SHPh) one and, thus, the formerly (very high) pressure falls down rather abruptly, the formerly cold nuclear medium starts imploding almost freely into the new vacuum. If the star mass is sufficiently large, then this implosion is shown to result in an enormous heating - up to the temperature about 100 MeV or, may be, even higher - and growth of the inner pressure due to degeneracy breaking and multiple $q\bar{q}$-pair production which withstands the gravitational compression (remind that the highest temperatures of supernovae bursts, as well as of the "normal" NS, are, at least, of one order lower). As a consequence, a "flaming wall" is, most probably, emerged on the way of further collapsing which prevents the NS to evolve towards the BH horizon appearance. At the same time, it could give rise to the most powerful GRBs produced by some very distant (young) stars.

1 Two incompatible mechanisms of neutron star instability

Two mechanisms underlying a compact star instability are confronted below which make the star to evolve in absolutely alternative ways: the first one implies HPh $\rightarrow$ SHPh transition within nuclear matter (it is described here in more detail) and the second one, which is rather familiar, is NS shutting to BH. Thus, the main point is to understand, which one is activated before.

1.1 Phase transition in nuclear medium

Schematically, this transition is depicted as follows:

$$
\begin{array}{ccc}
\text{QCD HPh} & \Longleftrightarrow & \text{QCD SHPh} \\
\\
\Downarrow & & \Downarrow \\
\\
P_{vac}^0 = -\varepsilon_{vac}^0 \simeq 5\,10^{-3}\ \text{GeV}^4 & \Longleftrightarrow & P_{vac} = -\varepsilon_{vac} \rightarrow 0 \\
\\
\Downarrow & & \Downarrow \\
\\
P_{tot} \simeq P_{vac}^0\ [rarefied\ gas\ of\ nucleons] & \Longleftrightarrow & P_{tot} = P_{vac}\ +\ P \\
\\
\varepsilon_{tot} \simeq \varepsilon_{vac}^0\ [rarefied\ gas\ of\ nucleons] & \Longleftrightarrow & \varepsilon_{tot} = \varepsilon_{vac}\ +\ \varepsilon
\end{array}
$$

Here $(\varepsilon_{vac}^0,\ P_{vac}^0)$ and $(\varepsilon_{vac},\ P_{vac})$ stand for the vacuum parameters (energy density, pressure) in HPh and SHPh, respectively, while $(\varepsilon,\ P)$ are the particle ones, and $(\varepsilon_{tot},\ P_{tot})$ are the overall energy density and pressure within the nuclear medium. It is worth emphasizing that $|\varepsilon_{vac}^0| \simeq \varepsilon_n$, the latter being the particle energy density of "close packed" nucleons (neutrons) which is somewhat - most probably, (25-30)% - higher than the mean intrinsic energy (mass) density within a free nucleon itself.

Below, we consider two conceivable scenarios of this phase transition [1, 2, 3] - the hard scenario, when the HPh transforms at some fixed density (pressure) directly (stepwise) into the current quark state (this is the "conventional" phase transition), and the soft one, which admits an intermediate state in between (a kind of crossover). The latter asks for the notion

of deconfined dynamical quarks (valons) - quasi-particles of non-fixed mass, which diminishes along with the density (pressure) increase. It is shown below that both scenarios result in developing strong instability under the phase transformation.

1. ***Hard scenario***:

    stepwise transition to the "empty" vacuum where $P_{vac} = \varepsilon_{vac} \equiv 0$

This scenario implies that the current quarks (almost massless (u, d)- and ~150-MeV s-quark) are emerged promptly as neutrons crush down. It is illustrated in Fig.1 [1] where the HPh vacuum condensate pressure P_{vac}^0 is confronted with that of the degenerate current quark gas at the different particle number densities, which is represented by the quark specific volume $\langle v \rangle$. It is obvious, that, in the framework of this scenario, the transition into *degenerate ("cold") quark gas* is ruled out - in fact, the matter starts collapsing into the new zero-rigidity ("empty") vacuum, what gives rise to an enormous heating (see an estimate below) of the nuclear medium just after the phase transition point is passed through [1].

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[1]Two pressures - the HPh-vacuum pressure and the pressure of degenerate perfect SHPh-particle gas - would equate only at the quark number density (point B in Fig.1), which is 3-4 times as high as the phase transition point one. It is worth noting that the neutrinos get essentially stuck under relevant densities and, thus, there is no way for an "instant" energy release from the star interior.
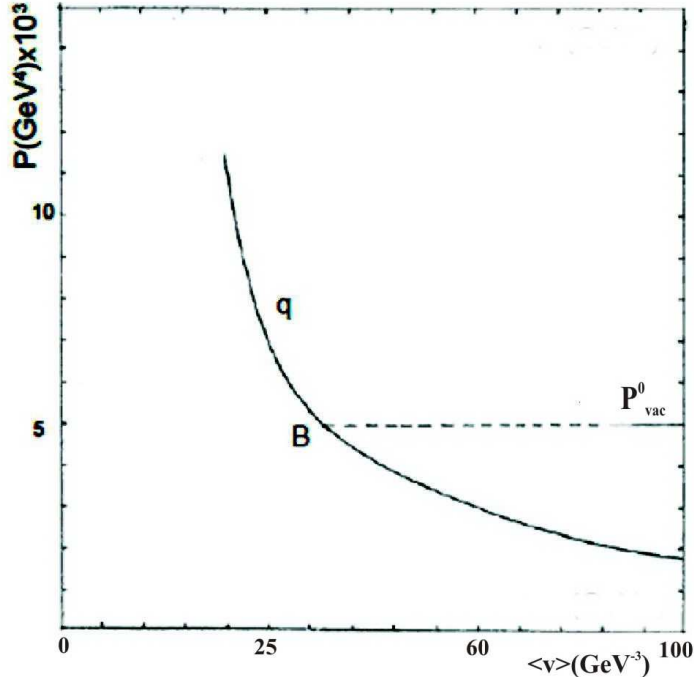
Figure 1: The pressure of the HPh non-perturbation QCD vacuum condensate (horizontal segment P^0_{vac}) vs the pressure of the degenerate ("cold") perfect gas of (u, d, s) current quarks (curve q). As the particle number density approaches the critical value (the neutron spacing becomes compact just as the quark specific volume is $\langle v \rangle \simeq 100$ GeV^{-3}), the occurrence of a gross gap between the HPh- and SHPh-phase pressures is unambiguously pronounced - at $\langle v \rangle \simeq 100$ GeV^{-3} the former is about three times as large as the latter one.

2. *Soft scenario*:   No stepwise HPh $\longleftrightarrow$ SHPh transition

This scenario implies that an intermediate state is to be passed after the neutrons "get in touch with each other". Namely, first, the neutrons disintegrate into the deconfined massive dynamical quarks (valons) [4, 5, 6, 7, 8] and, then, both the valon masses and vacuum condensate pressure decrease similarly along with the medium density increase [2]; finally, the valons turn into the current quarks and, thus, the vacuum condensate vanishes, $P_{vac} = -\varepsilon_{vac} \rightarrow 0$, more or less gradually.

A reasonable approach was suggested [2] which would describe the degenerate valonic gas at the particle energy densities $\varepsilon \geq |\varepsilon^0_{vac}|$ [2]. It is based on the EoS for a perfect gas with effective particle mass varying along with the medium density variation:

$$\varepsilon = \frac{6N_f}{2\pi^2} \int_0^{p_F} dp\, p^2 \sqrt{p^2 + m^2(\varepsilon)}, \tag{1}$$

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[2]If a steady state of this type were ever accessible - we shall argue below that, actually, it is not the case.

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where $N_f = 3$ is the number of flavors allowed for and the Fermi momentum $p_F = (\frac{\pi^2}{N_f \langle v \rangle})^{1/3}$. Let us parameterize the valon mass as

$$m_{u,d} \simeq m_0 \exp[-a\,(\varepsilon/|\varepsilon_{vac}^0| - 1)], \tag{2}$$

and, correspondingly,

$$\varepsilon_{vac} \equiv -P_{vac} \simeq \varepsilon_{vac}^0 \exp[-a\,(\varepsilon/|\varepsilon_{vac}^0| - 1)] \tag{3}$$

where $m_0 \simeq \frac{1}{3}m_n \simeq 330$ MeV [3] and a is a free parameter, which describes the rate of QCD vacuum condensate destruction. The numerical solutions of eq.(1), supplemented with eq's.(2,3) and the thermodynamic relation $P = -\partial(\langle v \rangle \varepsilon)/\partial \langle v \rangle$, are presented in Fig.2. Note, that only values $a \geq 1$ are physically reasonable because the HPh vacuum condensate should be significantly affected as the particle energy density approaches the absolute value of the condensate strength itself (or even earlier). The curves 2-4, which refer to $a < 1$, are depicted for an illustration only. It is evident that hard scenario comes back in the limit $a \to \infty$.

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[3]Actually, the $\sim$150-MeV mass-difference between (u, d)- and s-valons was allowed for, but no significant correction was shown to come therefrom [2].
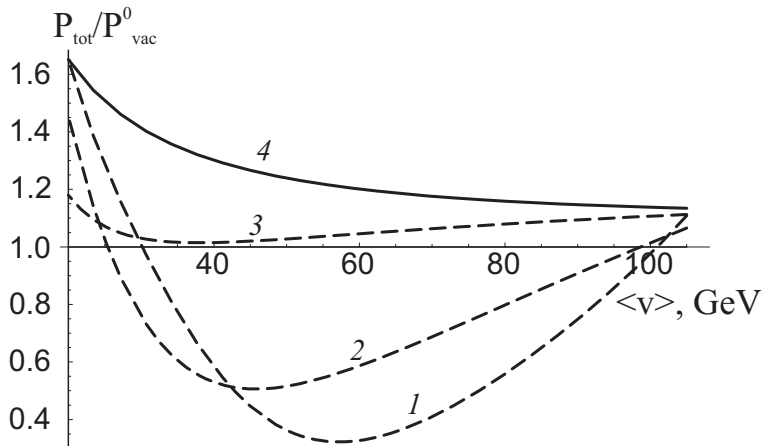
Figure 2: Formal solutions for the soft-scenario EoS at different rates of phase transition (HPh $\longrightarrow$ SHPh) in the "*cold*" (degenerate) nuclear matter. The curves (1 - 4) refer to the different regimes of the vacuum condensate destruction as well as of the valon mass decrease along with growing up the particle (valon) energy density ε (at $a = 1, 0.5, 0.1, 0.01$, respectively). One can see that a steady state regime of "*cold*" nuclear matter compression ($dP_{tot}/d\langle v \rangle < 0$) is quite improbable: it asks for unreasonably robust vacuum condensate, which would keep itself almost unchanged until the particle energy density becomes, at least, about one order higher than the module of HPh-condensate energy density ε_{vac}^0 ($a \leq 0.1$). Within the defensible framework, which implies that this condensate is subjected to significant destruction as both energy densities become of the same order (the downmost curve 1, $a = 1$) or even earlier, the inequality $dP_{tot}/d\langle v \rangle > 0$ would hold inevitably for a certain stage to be passed under the gravitational compression. This fact signals, undoubtedly, of instability. In other words, the nuclear substance can no longer remain "cold" (degenerate).

Thus, basing on the set of solutions shown in Fig.2, we come to the principally significant conclusion that *no way is conceivable for "cold" HPh $\longrightarrow$ SHPh transition* [4].

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[4]Observed EoS softening towards the center of large mass NS [9] could be considered as a certain phenomenological manifestation in favor of this statement, although the authors themselves suggest a different reasoning for this fact.

1.2 Blowing up NS *versus* BH formation

Nevertheless, one can expect that a certain transient "quasi-steady" state of NS (instead of the immediate NS rupturing) could be still maintained for some range of large M_{NS}, despite of enormous thermal disbalance between some small hot central domain and the "rest" of the star.

Is this pattern favorable for leaving NS a loophole to evolve towards the final BH configuration due to fast gravitational compression, which could enclose the horizon regardless of blowing up made by the divergent heat flow? Below, we try to put forward some meaningful arguments that the most reasonable answer must be negative.

As being emerged at the star center, the SHPh domain starts swelling until a transient balance is established between further heating due to gravitational compression and a pretty slow heat outward transport [5]. The inner (SHPh) domain is a kind of very hot subhadronic matter, which is called below, for brevity, QGP [6]. If the temperature of this plasma is T, then the obvious energy-conservation equation reads:

$$- AG\frac{M_{NS}^2}{R^2}dR \simeq 4\pi\sigma_{QGP}T^4\left(1 + \frac{|\varepsilon_{vac}^0| - \varepsilon_n}{\sigma_{QGP}T^4}\right)r^2 dr, \tag{4}$$

where on the left-hand side stands the work made by the gravitational field (M_{NS} and R are the NS mass and its radius, respectively, and the value of coefficient A is confined in between of its non-relativistic and ultra-relativistic limits, $\frac{6}{7} \leq A \leq \frac{3}{2}$ [7]), while on the right-hand side stands the energy increase within the domain of radius r occupied by SHPh (QGP),

$$\sigma_{QGP} = \frac{\pi^2}{30}\left(2 \times 8 + 2 \times 3 \times 2 \times 3 \times \frac{7}{8}\right)$$

being the 3-flavor QGP weight factor (8 gluons of spin 1 and $(3 + \bar{3})$ colored quarks of spin 1/2).

Nearly outside of the phase transition "boundary" (which is, actually, not a boundary but rather an extended spherical layer), the energy density of HPh substance made of closely

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[5]Note again that neutrinos get essentially stuck at the relevant densities of nuclear matter and, therefore, this transport is an extremely slow process (tens of hours *vs* the typical hydrodynamic-time scale, which is, probably, of some milliseconds).

[6]In the present context, it is referred to be a nearly perfect gas, which consists of the unremovable "primordial" quarks (the net baryon-over-antibaryon surplus) as well as of the multiply produced qluons and $q\bar{q}$-pairs, baryonic chemical potential μ_B thus tending to zero. Actually, the reasoning we put forward keeps valid for any microscopic structure, which mimics macroscopically the perfect gas thermodynamic behavior.

[7]Below, we put $A = 1$, since, in fact, the ultra-relativistic limit is rather inaccessible for the HPh-medium.

packed neutrons approaches the value $\varepsilon_n \simeq |\varepsilon_{vac}^0|$, what refers to $a \simeq 1$, in eq.s (2,3). This is just about the total density because the QCD vacuum condensate tends to zero. On the other hand, on the layer inside, it equals to the energy density of high-temperature QGP. Thus, the hydrodynamic (fast process) balance asks strongly for the leveling of two these energy densities:

$$|\varepsilon_{vac}^0| \simeq \sigma_{QGP}T^4, \tag{5}$$

wherefrom one obtains $T \simeq 130$ MeV [8], and we see that this balance could be only maintained at the price of an enormous thermal disbalance (remind that the NS medium temperature outside of the phase transition "boundary" is of a few MeV). It is worth emphasizing that this estimate is quite compatible with the $\mu_B = 0$ lattice MC simulation result [10] for HPh $\rightarrow$ SHPh transition, which was found to be a crossover lasting over the temperature interval $140\text{MeV} \leq T \leq 200\text{MeV}$.

Basing on eq.(5), one can reasonably assume that the second term in the brackets on the right-hand side of eq.(4) is small as compared to unity. If so, then the transient "quasi-steady" mode of a high-mass NS nuclear medium is composed as follows:

$$G\frac{M_{NS}^2}{R} =\simeq \frac{4\pi}{3}\sigma_{QGP}T^4 r^3 + C, \tag{6}$$

where C is defined by $\overline{M_{NS}}$ - the value of mass upper limit for the *really* stable ("cold", i.e., $r = 0$) NSs: $C \simeq (0.5 \div 1)\,M_\odot$ for $\overline{M_{NS}} \simeq (1.5 \div 2.5)M_\odot$ and $R \simeq (8 \div 10)$ km, respectively. Of course, the correlation (6) between M_{NS} and r can be defensible at $r \ll R$ only. In this case, a rather "peaceful" evolution of NS is not ruled out; it gives rise to the production of some cannonballs and/or successive GRBs, which become, however, the more destructive the larger is M_{NS}, thus resulting in diminishing the NS mass until it approaches the upper limit $\overline{M_{NS}}$ of NS stability. Otherwise (at still higher NS masses, when eq.(6) would predict r and R of the same order), no transient hydrodynamic balance is conceivable at all - the development of powerful shock waves seems inevitable which should forward NS towards the catastrophic self-destruction [9].

At the same time, the elementary condition for horizon first appearance within the body of a compact star reads: $\frac{2GM_g}{R_g} = 1$, or

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[8]This is, at least, one order higher than the typical temperatures for the supernovae explosions.

[9]From the more general point of view, all that is nothing else than different ways of symmetry (in this case - the chiral one) breaking along with the medium cooling: the no-order-parameter SHPh turns into the HPh, which shows up clearly an order parameter - it can be chosen to be the inverse radius of color confinement.

$$R_g \simeq [\frac{3}{8\pi G\langle\varepsilon_g\rangle}]^{1/2,} \tag{7}$$

where R_g and $\langle\varepsilon_g\rangle$ are the BH radius and its mean energy density, respectively. For getting the lower estimate of R_g, one has to take into account that $\langle\varepsilon_g\rangle \leq |\varepsilon_{vac}^0|$, since, otherwise, the phase transition instability followed by the aforementioned destructive cataclysms is expected to activate before. Thus, one obtains

$$R_g \geq 12 \text{ km or } M_g \geq 4\, M_\odot$$

We see that the admissible maximal NS and minimal BH radii are rather compatible, while the corresponding masses are separated by a significant gap. What is in between? If the NS mass were imagined to access $4\, M_\odot$, then eq.(6) tells immediately that $r \simeq R$, what is, actually, meaningless. Therefore, the only interpretation, which remains within eyeshot, is that the star can not "jump over" the gap without being ruined in full.

As for the stars of lower $\langle\varepsilon_g\rangle$ (larger M_g and R_g, both being $\sim \langle\varepsilon_g\rangle^{-1/2}$), the question remains open, since, in this case, a more detailed description of star dynamics should be involved. What can be said, is that shutting to BH seems to be, actually, an event of even lower probability than it might be thought of using the above consideration. The matter is that the horizon appearance is linked to the *global* features of the star nuclear medium (the value of M/R and *averaged* energy density $\langle\varepsilon\rangle$), while the HPh $\longrightarrow$ SHPh transition instability is linked directly to the *local* value of ε, which is, undoubtedly, r-dependent and increases towards the star center. The same argument gives, all the more, "obvious preference" to the proactive development of HPh $\longrightarrow$ SHPh instability in case of some density fluctuations within the star body. Thus, this instability is anyway expected to start developing at lower values of $\langle\varepsilon\rangle$.

It is worth also mentioning, in this connection, that a number of additional factors - possible star non-sphericity, star rotation and, especially, binary-star configuration - should, obviously, result in diminishing the margin of NS stability, thus making the above arguments against a BH-horizon appearance even more defensible.

2 Conclusion

The QCD-induced mechanism of additional NS instability is discussed. The NSs of highest masses are proven to be in face of instability associated with QCD-vacuum transformation under HPh $\longrightarrow$ SHPh transition, which could manifest itself, in particular, through the soft-

ening of EoS towards the star center. This instability seems to develop before a BH horizon appears within the star body, what makes rather improbable the very accessibility of a BH configuration at the end of collapsing star evolution.

Since the SHPh-phase temperature is, at least, one order higher than that of the super-novae explosions, the expected energy release could be several (seemingly, up to 2-3) order higher. That is why it is difficult to resist the temptation of linking the QCD-induced insta-bility under discussion and the poorly understood data on very distant ("young") GRB's of highest energy, like GRB 090423 [11], GRB 080916C [12], GRB 080319B ("naked eye") [13], etc.

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