We summarize recent developments in identifying the ground state of dense baryonic matter and beyond. The topics include deconfinement from baryonic matter to quark matter, a diquark mixture, topological effect coupled with chirality and density, and inhomogeneous chiral condensates.

Keywords: High baryon density, Nuclear matter, Quark matter, Topological effects

1. Onset of deconfinement at high temperature

Color confinement in QCD is still a big mystery and some people may want to think that color deconfinement would be more understandable; otherwise, we have no idea what we are talking about when it comes to the so-called quark-gluon plasma (QGP). It is our belief that the QGP has been created in the laboratory by means of nucleus-nucleus collision at high enough energy at RHIC in Brookhaven and LHC in CERN. It is not our present aim to challenge this widely accepted interpretation. The goal of this contribution should be to encourage for physics opportunities of the lower-energy nucleus-nucleus collision mainly in the context of the QCD phase diagram research. Since our theoretical understanding in finite-density QCD is so poor, what we can do the best is to make a guess from what we have known. In other words, we must carry extrapolation out from known physical states, or preferably interpolation between known states if possible.

In this section we discuss the extrapolation from high temperature and zero density along the phase transition. Nowadays it is said that the lattice QCD simulation has unveiled all features of QCD thermodynamics as long as the baryon chemical potential is much smaller than the temperature $T$. Indeed the lattice QCD simulation is so powerful that precise data of the pressure, the internal energy, the entropy density, etc are available as functions of $T$. Historically speaking, because the QCD transition temperature turned out to be $T_c \sim \Lambda_{\text{QCD}}$ in early lattice QCD studies, the dream for the QGP factory came to reality [1, 2]. Figure 1 is a schematic figure of the dimensionless pressure $p/T^4$ that effectively measures the physical degrees of freedom in the thermal system. It is quite conceivable to interpret the rapid rise in $p/T^4$ as the liberation of colored particles; $N_c$ quarks and $(N_c^2 - 1)$ gluons.

Thus, answering the following question seems easy; where do you find color deconfinement? Your answer would be that it takes place at $T_c \sim \Lambda_{\text{QCD}}$ where $p/T^4$ starts growing up. This conservative answer is of course not wrong, but not completely satisfactory. For the increasing behavior of $p/T^4$ another interpretation is possible, as illustrated by the HRG line in Fig. 1(a) that schematically describes the thermodynamics predicted in the hadron resonance gas (HRG) model. The HRG assumes a gas of non-interacting hadrons in a thermal bath. Each hadron hardly contributes to the entire thermodynamics, but the whole sum from hundreds of hadrons amounts to a substantial portion of the total pressure. As a side remark I note that it is a bit subtle whether the hadron interaction is incorporated or not in the HRG model. Some people claim that interaction is correctly taken into account via higher resonating states.
Such a claim is true if all the $S$-matrix poles are picked up by existing bound states and/or resonances. There may be, however, missing contribution from branch cut associated with threshold behavior and from hidden states not listed on the particle database. The theoretical foundation for the validity of the HRG estimate should deserve future investigations.

In any case, the fact is that the coincidence between the lattice QCD and the HRG is becoming better and better up to $T_c$ as the lattice QCD approaches the continuum limit and the physical quark mass. Above $T_c$, eventually, the HRG badly blows up, which is to be regarded as the breakdown of the HRG model. Here, I would point out two non-trivial problems in this HRG results, which will provide us with a useful insight when we address dense matter.

The first problem is that we can no longer consider that the rapid rise in $p/T^4$ should be attributed to color deconfinement. It could be induced by hundreds of hadrons even without colored particles at all. This is why I wrote; it is not wrong but not completely satisfactory to say that deconfinement takes place when $p/T^4$ shows a quick increase. Then, you may wonder, as a radical extreme, if it makes logical sense that the system keeps hadrons as relevant degrees of freedom for any $T$ including $T > T_c$. We already know that the standard HRG is not valid above $T_c$, but the question is what exactly breaks down there. Does the hadronic description lose its meaning at all or is it only a part of assumption made in the HRG that goes invalid?

This question is deeply related to my second problem; that is, how is it possible that the thermodynamics is taken over from the HRG to the QGP as observed in the lattice QCD [see Fig. 1(a)] even though the pressure of the HRG is greater. It is a firm thermodynamic principle that any state with a larger pressure would be more favored. Therefore, if the HRG is somehow legitimate above $T_c$, the hadronic state should be more stable than the QGP, which is not the case in reality. It is not very difficult to propose a consistent picture that resolves my two problems. Hadrons or mostly mesons are accompanied by interaction clouds inside of a non-zero radius $\sim \Lambda_{\text{QCD}}^{-1}$. Thus, the Hagedorn-like behavior of the HRG pressure should be saturated by the interaction clouds of mesons. It follows that each border of meson is not really well-defined and quarks and gluons can hop around between overlapping wave-functions of blurred mesons. This is nothing but a percolation picture of color deconfinement. So, we do not have to abandon the hadronic description even above $T_c$ only if we incorporate not point-like mesons but extending mesonic wave-functions, and this is equivalent to introducing colored particles in the thermal system after all.

Now we have a right answer to the question about the onset of deconfinement. Your answer should be that deconfinement takes place when the growth in $p/T^4$ is saturated due to hadronic interactions. Such a strongly interacting mesonic state is more naturally and efficiently handled by (resummed) perturbation theory of hot QCD and the hard thermal loop approximation is so successful to account for the lattice thermodynamics from the high temperature side up to the temperature around $\sim 2T_c$ as sketched in Fig. 1(b). We notice that the most non-trivial region near $T_c$ is to be characterized by a sort of duality between strongly interacting mesons and weekly interacting quarks and gluons. Because this region looks so special, we are tempted to name it, and actually a well-known name has already been given; that is, the strongly-correlation QGP or shortened as the sQGP. Some other people named it as the semi-QGP and...
this latter is much better than calling it the sQGP in my opinion. Here, I would propose another name; the Gluesonic Matter in a sense that gluonic degrees freedom $\sim O(N_c^2)$ can be realized in the interacting mesonic matter. You may think that it is just a matter of terminology, but I would stress that many people still misunderstand the essential idea of the Quarkyonic Matter, as we will discuss later, and the alternative name of the Gluesonic Matter helps us with understanding the Quarkyonic Matter correctly.

2. Onset of deconfinement at high baryon density

We are ready to go into our main discussions on dense QCD matter. Quark matter was first speculated in neutron stars [4] simply as a hypothetical state of matter. You may want to argue that the asymptotic freedom at high density and the Debye color screening could justify deconfined quarks and gluons like the QGP at high temperature [5]. The contemporary QCD theorists are aware, however, that magnetic gluons cannot be screened and it may be even misleading to identify the typical scale of interaction with the chemical potential ($\mu_B$ for the baryon number or $\mu_q = \mu_B/N_c$ for the quark number).

It is thus a profound and unanswered question how we can retrieve any information on density-induced deconfinement from QCD itself. Theoretically speaking, even if the lattice QCD were at work to study finite-density systems, the Polyakov loop that is an approximate order parameter for deconfinement at high temperature would be of no use in order to characterize a state of matter in the low-$T$ and high-$\mu_B$ region. Then, how can we see deconfinement with any physical observable?

Let us recall that we have witnessed deconfinement not relying on an approximate order parameter in the previous section. That is, deconfinement from baryonic to quark matter should take place when the pressure is saturated to be of $O(N_c^2)$ due to baryonic interactions. Interestingly enough, the large-$N_c$ expansion of QCD rigorously justifies such a physics picture.

In the large-$N_c$ limit baryons interact strongly via quark exchange. The constraint that baryons carry no net color allows for $N_c$ combinations of quark exchange (or $N_c^2$ combinations with one gluon exchange leading to $g^2 \sim 1/N_c$) as seen in Fig. 2(a), which indicates that the baryon interaction energy is of $O(N_c)$. Now it is obvious that the situation is quite analogous to what we have seen using the HRG model at low-$\mu_B$ and high-$T$. Therefore, baryonic matter is a dual of quark matter as it is, and in this sense, we should reasonably call it the Quarkyonic Matter [6]. If you do not like the name of the Quarkyonic Matter, you can give your own name such as the “strongly-correlated quark matter” etc, but it is just a matter of terminology.

Apart from chiral symmetry that is the subject in later discussions, we can draw a qualitative QCD phase diagram as depicted in Fig. 3. The hadronic matter is surrounded by the chemical freeze-out curve that is associated with the Gluesonic Matter in the low-$\mu_B$ and high-$T$ region, the Quarkyonic Matter in the low-$T$ and high-$\mu_B$ region, and a transitional region between them that looks like a triple point or a Triple Region [7].
There is one important distinction of the Quarkyonic Matter from the Gluesonic Matter. In the large-$N_c$ limit mesons are non-interacting particles because the decay constant scales as $f_\pi \sim \sqrt{N_c}$ and so the interaction vertices should be suppressed accordingly to make the total amplitude stay finite. This is why the pressure has a jump as a function of $T$. As long as the mesons are excited dilutely, the interaction clouds are negligible and the pressure remains to be of $O(N_c^0)$. In contrast to this, in the Quarkyonic Matter, baryons always interact strongly. So, it is quite unlikely that a rapid rise as in Fig. 1(a) occurs with increasing density or chemical potential. We should therefore consider that the Quarkyonic Matter extends smoothly from conventional nuclear matter to quark matter at asymptotically high density (and this is why some people correctly say that the Quarkyonic Matter is nothing but nuclear matter, which is not wrong but not completely satisfactory). I would emphasize that this is a good news for future experimental prospects. We were excited about the so-called sQGP and why not should we about the counterpart in the high density region that has an even wider terrain on the phase diagram? It is not easy to make any concrete predictions on the properties of the Quarkyonic Matter, but we should remember that no concrete predictions had existed before RHIC that would signal for the sQGP or the Gluesonic Matter.

Coming back to the theoretical consideration, those who are familiar with dense QCD may have the following question; where is the color superconductivity (CSC)? This is an absolutely decent question. I know that some people living in the strict large-$N_c$ world would emphasize the irrelevance of CSC too much, making too little of the reality of CSC. I believe, however, that we are living in a world with $N_c = 3$ and we must directly face such a question of how to reconcile the Quarkyonic Matter and the CSC on equal footing.

3. Reality of diquarks

The picture of the Quarkyonic Matter is not necessarily incompatible with the presence of diquarks. Since the recognition of the Quarkyonic Matter originated from the large-$N_c$ limit where the diquark interaction is suppressed, you may well think that the Quarkyonic Matter and the CSC are not cooperative but behave more like cats and dogs. Of course, cats and dogs would not always struggle with each other, and there should be a way how they could live together peacefully beneath the same roof.

In the large-$N_c$ world the sources of the baryon density and the pressure are clearly separated; the baryon density appears solely from static baryons and the pressure is dominated by the mesons between baryons. Although it is not written explicitly in Fig. 2(a), multiple gluons propagate between exchanged quarks in the $t$-channel, and the interaction is mediated by not two quarks but a meson. So, our intuition based on a Fermi gas with Fermi surface completely breaks down. From the diagrammatic point of view, however, there is no reason why we should exclude a resonating state of two quarks in the $s$-channel. Such an intermediate state of two quarks is not a color singlet, but as long as it emerges through a virtual state, there is no practical problem. Then, with ladder-type resummation of gluonic processes, Fig. 2(a) actually describes:

- Two baryons having a strong interaction mediated by meson exchange in the $t$-channel.
Two baryons having a mixture with two \((N_c - 1)\)-quark objects and a diquark in the \(s\)-channel.

Here you may concern that the four-quark interaction in the diquark channel should be suppressed by \(1/N_c\) than that in the meson channel, but this \(1/N_c\) is compensated for by \(N_c\)-colored diquarks. We see that the latter picture of a diquark mixture is more understandable on the intuitive level. In the \(N_c = 3\) world, as illustrated in Fig. 2(b), the interpretation is even more intuitively appealing. The \((N_c - 1)\)-quark part is also the diquark and a quark-diquark mixture represents the interacting intermediate state. Hence, we shall adopt a working definition of the Quarkyonic Matter characterized by a mixture of diquarks.

It is an empirically established idea to construct the baryon wave-function as a bound state of a diquark and a quark. In the quark model such a construction simply refers to the group theoretical structure of color indices, and it does not necessarily require the reality of diquarks. In the Faddeev equation in the four-quark interacting model the diquarks acquire more reality along the line of Fig. 2(b). Besides, the famous \(\Delta I = 1/2\) rule of non-leptonic weak decays suggests the presence of spatially compact diquark inside of the baryon wave-function.

Once we admit the presence of strong diquark correlation, we have an immediate problem. If we can define the diquark mass (or precisely speaking, if the diquark spectral function has a prominent peak at some frequency), and if \(\mu_q\) exceeds roughly \(N_c/2\) times the diquark mass, we cannot avoid the Bose-Einstein condensation of diquarks. This means, if the diquark mass is not so far from \(2/N_c\) times the baryon mass or twice of the constituent quark mass, the CSC is unavoidable even in the vicinity of nuclear matter. At a first glance you may be inclined to think that the CSC near nuclear matter is an artifact of the naive diquark model. Before making any judgment, we should think twice about the reality of the CSC in nuclear matter, however.

It is very unlikely that nuclear matter exhibits a sharp first-order phase transition to quark matter, though it was a conventional approach to postulate the equation of state of dense matter \[8\]. If we introduce the diquark degrees of freedom between nuclear and quark matter, everything seems to be quite consistent. As we already discussed, deconfinement cannot be usually guaranteed by the screening phenomenon in quark matter, but if the CSC (or the color-flavor-locked state, strictly speaking) occurs, all gluons are gapped as a result of the Meissner effect. Then, the perturbation theory becomes well-defined and the magnetic sector leads to the celebrated enhancement of the gap energy. In this way, it is not such a surprising proposition to define deconfinement of dense matter by the formation of the diquark condensate. Then, a CSC expert would pose the following question; how can you define the CSC or the diquark condensate in a gauge invariant way? This is not an academic question but rather a very pragmatic question. On the academic level the simplest answer is that I fix a gauge so that the diquark condensate can take a non-zero value. In real experiment, however, the physical observable should be gauge invariant and the diquark condensate is not detectable in principle. This undetectability is nicely summarized in terms of the Quark-Hadron Continuity \[9, 10\]. From this theoretical point of view the Quark-Hadron Continuity, in fact, there is no contradiction even if we assume a small fraction of the diquark condensate inside of ordinary nuclear matter. In nuclear theory one of the most well-known calculation schemes is the Hartree-Fock-Bogoliubov theory and the pairing interaction leads to a finite pairing gap. The diquark condensate would break the U(1) symmetry associated with the baryon number conservation as well as chiral symmetry. Because we know that both the U(1) symmetry and chiral symmetry are broken in superfluid nuclear matter, no symmetry argument prohibits the existence of the diquark condensate. In other words, if the existence is not ruled out by the symmetry reason, we should think that it must be there.

The reality of diquarks in nuclear matter would open an intriguing opportunity for upcoming experimental attempts to pursue compressed baryonic matter. Diquarks are fundamental building blocks of exotic hadrons composed from more constituent quarks than \(qq\) or \(qqq\). There is an idea to manifest the \(qq\) part only by inserting a heavy flavor \(Q\) into a state; \(Qqq\), which is actively promoted by Japanese hadron physicists (learned from private communications with A. Hosaka and M. Oka). The high baryon density is another way to manifest the diquark correlation. Here, it must be mentioned that the diquark correlation is usually seen in momentum space, as in the Cooper pair of electrons in the ordinary BCS theory. Hence, whenever we talk about the experimental challenge to see the diquark correlation, it should be clearly stated which of the correlation in momentum space and in configuration space is sensitive to the proposed observable. We should always keep in mind that demanding the presence of spatially compact diquark is a very strong assumption and there is no theoretical argument that can endorse any strong diquark correlation in space, that is the case also when the CSC is turned on. Nevertheless, the diquark should be enumerated as the top-priority keyword for the expected paradigm shift from the Gluesonic Matter to the Quarkyonic Matter in the near future.
4. Topological effects coupled to the baryon density

Let us change the subject from the state of matter to a fancy phenomenon with the quantum anomaly. In condensed matter physics the Magneto-electric Effect has been long known; if you impose $E$ on a special material, you observe $B$ in parallel to $E$, which implies a source term, $\theta E \cdot B$, in the effective Lagrangian. Such a term is called the $\theta$-term and the topological insulator corresponds to the case with $\theta = \pi$, for example. In the nucleus-nucleus collision, a very strong magnetic field whose energy scale is comparable with $\Lambda_{\text{QCD}}$ is expected if the impact parameter is non-zero. Then, if the QGP accommodates the $\theta$-term, it is quite conceivable to anticipate some topological effect in analogy to the Magneto-electric Effect.

The most well-known example of the topological effect along this line is the Chiral Magnetic Effect [11]. The severest problem in the Chiral Magnetic Effect is that it is sensitive to the Local Parity Violation and thus the net effect averaged over space and/or collision events should be vanishing. So, what we can see experimentally is only the fluctuation that is parity even, and in principle, the signal is not separable from the background. Even though the theoretical and experimental efforts are continuing for the Chiral Magnetic Effect, we further need a breakthrough to invent some better observable than the fluctuation of the charge separation; otherwise, we can never conclude anything for or against the Chiral Magnetic Effect.

The situation is significantly improved as soon as the density comes into the game of the quantum anomaly. The reason for this is obvious for theorists: the chemical potential is the zeroth component of the gauge field, and so, the finite-density physics is always accompanied by the gauge dynamics. A typical example is found in the density origin in (1+1)-dimensional gauge theories, in which a finite density arises purely from the Wess-Zumino term. Then, what is the theoretical prediction if we have a finite $\mu_B$ and a strong $B$? The answer is recently referred to as the Chiral Separation Effect and the axial current $j_5 \propto \mu_B B$ is generated. You may think that $j_5$ in the Chiral Separation Effect would induce the chirality imbalance in the same way as the charge separation in the Chiral Magnetic Effect, but the chirality is not a conserved charge. It decays through a finite mass and a topologically winding configuration. Still, the Chiral Separation Effect at finite density is much more advantageous than the Chiral Magnetic Effect because it is a net effect and the spatial and/or event average does not wash it out. As we stated, chirality is not a conserved charge, but the helicity is a good quantum number, and so the chirality decay may lead to an interesting consequence about, so to speak, helicity transmutation. Such a possibility has been addressed in the context of neutron star physics [12], and in my belief, this idea could be usefully imported to the analysis of lower energy heavy-ion collisions.

The physical interpretation of the topological currents has a subtle aspect. In theory the current just means an expectation value of spinor bilinear operator in a certain channel. If we have a condensate of the axial vector meson for some reason, we also acquire a non-zero $j_5$, and is this really a current that flows in real time? It is not so easy to let our imagination work, but the answer is yes. The vector condensation and the supercurrent are usually identifiable. Then, when we consider the ground state of dense QCD matter, the condensation of $j_5$ or its mean-field effect must be taken into account just like the mean-field density effect from the vector interaction. To this end, we have to move to our final topic; the possibility of the inhomogeneous chiral condensates.

5. Inhomogeneous chiral condensates

It is an old but still vital idea that the spatial modulation occurs at high density, so that the genuine ground state should be characterized by inhomogeneous chiral condensates. A nice review by pioneering researchers, Buballa and Carignano, is quite recommendable for interested readers [13]. The idea is traced back to the $p$-wave pion condensate in nuclear matter. Although the pion condensate is disfavored by the Gamow-Teller giant resonance, it is not yet completely excluded, and moreover, the idea is still alive as a realistic possibility inside of quark matter.

The simplest way to introduce spatial inhomogeneity is to postulate a one-dimensional spiral structure in the scalar and the pseudo-scalar channels as: $\sigma \sim \cos(qz)$ and $\pi^0 \sim \sin(qz)$ along the $z$ direction. Such a system is not much different from the homogeneous case, and in fact, if the density is calculated, it turns out to be a spatial constant. This type of chiral condensate is called the dual chiral density wave [14], or the chiral spiral. In the space of $\sigma$ and $\pi^0$ condensates, the chiral spiral looks like pasta called fusilli as in Fig. 4[1] I would say that the QCD ground state could be as fascinating as pasta al dente.

---

1In the talk at QM2015 I showed a famous painting by Escher, “Spirals”, which is a good analogy to the chiral spiral of QCD, and this beautiful
Here we are not going into details, but let us mention on the “facts” only. Some model calculations predict the so-called QCD critical point, and some others predict a smooth crossover on the whole QCD phase diagram without the critical point. This has been a status of the QCD critical point search, and unlike this frustrating situation, all the model calculations predict spatially inhomogeneous chiral condensates in the high-density region, which approximately overlays the region of the Quarkyonic Matter in Fig. 3. Such coincidence of the inhomogeneous region and the Quarkyonic Matter is not accidental; the chiral condensate in the large-$N_c$ world should be inhomogeneous. This can be understood from the fact that baryons are infinitely heavy in the large-$N_c$ limit. The optimal configuration of baryonic matter is thus a crystal of static baryons and then the chiral condensate becomes smaller locally near baryons than the vacuum value.

A frequently asked question about the inhomogeneous chiral condensates is that it may be unstable against mesonic fluctuations. This is a very reasonable question. In condensed matter physics, a hypothetical crystalline superconducting state (called the FFLO state) has been a long-standing issue. In most cases, such an exotic state has the lowest energy in the mean-field analysis, while it goes away once fluctuations of Nambu-Goldstone bosons or spatial rotations are considered. In QCD, therefore, we should pay a serious attention to the roles played by pion fluctuations and rotations of a finite-sized fireball to test the stability of the inhomogeneous chiral condensates. There are some works on this matter, but the final sentence still awaits to be announced.

Naturally, such a spiral would be affected by the topological current $j_5$ if $j_5 \neq 0$, since the chiral spiral represents a flow of chirality as perceived from Fig. 4. Therefore, when we talk about the chance to study baryonic matter in the heavy-ion collision, we should deal with both spatial modulations and topological currents. There are not many theoretical efforts yet in this direction, probably because of uncontrollable ambiguity in model treatments. So far, there is one concrete result in the large-$N_c$ limit using the holographic QCD model \cite{15}, which suggests that the spatial modulation becomes less favored with stronger axial-vector interaction, $j_5 \cdot j_5$, that should be significantly enhanced by the topological current $j_5 \neq 0$. This conclusion is consistent with the observation of the $p$-wave pion condensate diminished by the Landau-Migdal interaction in the spin-isospin channel.

If the inhomogeneous chiral condensates can survive in reality, the experimental survey for this structure should go in a similar fashion to the Local Parity Violation. Instead of parity-odd domains, we should search for dense bubbles. The problem is again that we cannot make a discovery in a qualitative sense but the analysis always relies on quantitative comparisons, and such a strategy does not work unless we establish undoubted theoretical predictions. This may sound like a difficult task, but I am rather optimistic. Before the age of RHIC and LHC, who could have foreseen such a big success of the thermal fit and the HRG model? We certainly need some refinement of those descriptions, probably with several mean-fields as in the relativistic mean-field model of nuclear matter or with new degrees of freedom like diquarks. With sufficient data of hadron multiplicity and their fluctuations at lower energies, it is simply a matter of time to come by a reliable baseline in order to diagnose the intrinsic properties of dense QCD matter. I would not guarantee anything particularly interesting between cold QGP and hot nuclear matter, that is the regime accessible by future heavy-ion programs, but I can at least say that this experimentally accessible regime could definitely provide us with lots of hints to physics questions that nobody can answer at present.

painting was supposed to be presented here. The M.C. Escher Company, however, charged 75 euros for copyright-fee, 35 euros for handling, and 75 euros additionally per electronic image. So, I gave up demonstrating artistic nature of physics and replaced it with another branch of human culture; gastronomic creation of Italy.
Lastly, I would like to express my thanks to G. Baym, M. Buballa, T. Hatsuda, A. Hosaka, L. McLerran, M. Oka, J. Pawlowski, M. Stephanov, T. Tatsumi, N. Yamamoto. They often gave me a hard time with unanswerable questions, sometimes opened my eyes to a new idea, and occasionally shared physics motivation with me. The contents of this contribution are largely influenced by stimulating discussions with these brilliant researchers.

References

[1] G. Baym, RHIC: From dreams to beams in two decades, Nucl. Phys. A698 (2002) XXIII–XXXII. arXiv:hep-ph/0104138 doi:10.1016/S0375-9474(01)01342-2
[2] K. Fukushima, Chiral Symmetry and Heavy-Ion Collisions, J. Phys. G35 (2008) 104020. arXiv:0806.0292 doi:10.1088/0954-3899/35/10/104020
[3] Y. Hidaka, R. D. Pisarski, Suppression of the Shear Viscosity in a "semi" Quark Gluon Plasma, Phys. Rev. D78 (2008) 071501. arXiv:0803.0463 doi:10.1103/PhysRevD.78.071501
[4] N. Itoh, Hydrostatic Equilibrium of Hypothetical Quark Stars, Prog. Theor. Phys. 44 (1970) 291. doi:10.1143/PTP.44.291
[5] J. C. Collins, M. Perry, Superdense Matter: Neutrons Or Asymptotically Free Quarks?, Phys. Rev. Lett. 34 (1975) 1353. doi:10.1103/PhysRevLett.34.1353
[6] L. McLerran, R. D. Pisarski, Phases of cold, dense quarks at large N(c), Nucl. Phys. A796 (2007) 83–100. arXiv:0706.2191 doi:10.1016/j.nuclphysa.2007.08.013
[7] A. Andronic, D. Blaschke, P. Braun-Munzinger, J. Cleymans, K. Fukushima, L. McLerran, H. Oeschler, R. Pisarski, K. Redlich, C. Sasaki, H. Satz, J. Stachel, Hadron Production in Ultra-relativistic Nuclear Collisions: Quarkyonic Matter and a Triple Point in the Phase Diagram of QCD, Nucl. Phys. A837 (2010) 65–86. arXiv:0911.4806 doi:10.1016/j.nuclphysa.2010.02.005
[8] G. Baym, S. Chin, Can a Neutron Star Be a Giant MIT Bag?, Phys. Lett. B62 (1976) 241–244. doi:10.1016/0370-2693(76)90517-7
[9] T. Schaefer, F. Wilczek, Continuity of quark and hadron matter, Phys. Rev. Lett. 82 (1999) 3956–3959. arXiv:hep-ph/9811473 doi:10.1103/PhysRevLett.82.3956
[10] M. G. Alford, J. Berges, K. Rajagopal, Unlocking color and flavor in superconducting strange quark matter, Nucl. Phys. B558 (1999) 219–242. arXiv:hep-ph/9903502 doi:10.1016/S0550-3213(99)00410-1
[11] K. Fukushima, Views of the Chiral Magnetic Effect, Lect. Notes Phys. 871 (2013) 241–259. arXiv:1209.5064 doi:10.1007/978-3-642-37305-3_9
[12] A. Ohnishi, N. Yamamoto, Magnetars and the Chiral Plasma Instabilities, arXiv:1402.4760
[13] M. Buballa, S. Carignano, Inhomogeneous chiral condensates, arXiv:1406.1367
[14] E. Nakano, T. Tatsumi, Chiral symmetry and density wave in quark matter, Phys. Rev. D71 (2005) 114006. arXiv:hep-ph/0411350 doi:10.1103/PhysRevD.71.114006
[15] K. Fukushima, P. Morales, Spatial modulation and topological current in holographic QCD matter, Phys. Rev. Lett. 111 (2013) 051601. arXiv:1305.4115 doi:10.1103/PhysRevLett.111.051601