

Chapter 1

Introduction

The core of this volume, the problem of “mass” of what we see around us, touches on one of the most profound issues in all areas of physics, be they condensed matter, nuclear or particle physics. This remains just as much of a mystery in nature as where the dark energy of the Universe comes from and what the dark matter of the Universe is. Though manifested differently in different areas of physics, it presents a basic puzzle. The origin of the mass of the particles familiar to nuclear physicists, namely, proton and neutron, belongs to this class of puzzles. Unlike the particles studied in high-energy machines, the proton is a mundane object, being a basic constituent of nuclei, that is accessible to laboratory scrutiny. Nonetheless its structure has remained up to date prominently mysterious. While the masses of all “visible” matters around us, molecules, atoms, nuclei *etc.* are more than 99% accountable in terms of the sum of the masses of the constituents, the nucleon being the basic constituent of them, the source of the nucleon (proton and neutron) mass itself is still more or less unknown. QCD tells us that a proton, say, is made primarily up of two up (u) quarks and one down (d) quark, with the strange (s) quark and heavier quarks entering into its structure only negligibly. Since the quarks involved have tiny masses, nearly zero on hadronic scale, $\sim 98\%$ of the proton mass must be coming from something other than the *basic* constituents of the proton. So the question is: Where do $\sim 98\%$ of the proton mass come from?

In CND-I, it was explained how the rearrangement of the complex vacuum can be responsible for the generation of the mass and for much of the properties of strong interaction dynamics as we observe them. The phenomenon was attributed to “spontaneous” break-down of chiral symmetry – which is an exact symmetry if one ignores the (tiny) quark masses. Although it is not yet proven, the dynamical breaking of chiral symmetry may also be intricately connected with the confinement of quarks and gluons. In this volume, we shall accept this possibility and ask, if the mass of hadrons is dynamically generated from the complex vacuum structure, can one unravel the mechanism of the mass generation by reversing the process, namely tweaking the vacuum so that the mass be made to dynamically “disappear”? One may think of this as going backwards in time in the evolution of the Universe from high temperature/density to what it is now.

In short, the main subjects of this volume consist of exploring how this can be done in the terrestrial and space laboratories and how to explain what have been so far observed in consistency in a language that models the fundamental theory of strong interactions.

In contrast to the time of the writing of CND-I when little empirical information was available, some, though incomplete, experimental data have accumulated which could now enable us to answer some part of the question posed above. Hadrons are being heated in heavy-ion collisions to high temperature, say, several hundreds of MeV, allowing us to have a glimpse of the Universe at the age of $\sim 10^{-6}$ second after the Big Bang. Hadronic matter is, and will be soon, compressed to a density several times the nuclear matter density, mimicking the conditions that are supposed to be present in the interior of compact stars. These processes bring hadronic matter to a temperature or density at which the spontaneously broken chiral symmetry is restored or/and the quarks and gluons get deconfined. Along the way the matter experiences a variety of conditions in which the mass and coupling constants get modified significantly, influencing crucially the properties of hadrons living in these environments. QCD should be undoubtedly the right theory to address all these issues but a direct use of QCD for *all* such issues is out of reach at this point. We are thus forced to resort to effective field theory (EFT for short) formalism.

We develop the appropriate languages to address this issue in this volume. We do this chiefly in two different ways. One way is to pick up what was touched on in CND-I, namely approaching effective field theory from the quark-gluon picture via what is called the Cheshire Cat Principle (Chapter 3). This is a “bottom-up” approach in the sense that we go up in energy scale to probe strong interactions at the relevant kinematic regime. The other is via holographic dual QCD based on AdS/CFT from string theory. This approach, touched on only briefly in Chapters 5 and 6, which is being actively developed at the time of writing of this volume may be classified as a “top-down” approach since it is based on a string-theory idea. The end product in both cases is effective chiral Lagrangian theory. In our case, the EFT tool to use is the hidden local symmetry (HLS) theory. We make a detailed exposition of the latter approach in Chapter 5. This may not be the only way to approach the issue, but we find this to be the easiest and most efficient way.

To make contact with the traditional nuclear physics, we treat few-nucleon nuclei in Chapter 4. Our theme is *continuity* in physics between the “old” way of doing things and the “new” way that is anchored on more modern view. In addition we aim to illustrate the power of the EFT approach in nuclear physics *proper* where the temperature is low and the density is not high. The focus will be on the predictive power of the EFT in nuclear physics. Here certain predictions – free of arbitrary parameters – are made by means of effective field theory combined with traditional many-body approaches based on accurate and sophisticated phenomenological techniques. The import of this exercise is that EFT can supplement and improve on – not replace – the standard nuclear physics approach (SNPA for short). What makes

EFT valuable in nuclear physics is that it can make predictions that go beyond what was feasible in standard nuclear physics approaches.

The skyrmion picture that figured centrally in CND-I reappears in a different guise in Chapters 6, in particular in making certain predictions of novel structure at high density. The skyrmion as a baryon re-emerges in the gravity dual description of QCD that comes from string theory. What makes this development noteworthy is that in a realistic holographic dual QCD (or AdS/QCD), low-energy strongly-coupled dynamics is captured by hidden local symmetry theory with an infinite tower of vector mesons coupled to the Goldstone pions (HLS $_{K=\infty}$ for short) and the entire baryon degrees of freedom are encoded in the theory as solitons. The soliton in that theory is an instanton in 5D Yang-Mills theory or equivalently a skyrmion in 4D HLS $_{K=\infty}$ theory. It turns out that the effect of infinite tower is totally encapsulated in a 5D instanton and gives rise to a baryon structure which is *drastically different* from that of the Skyrme model based on pion-only chiral Lagrangians. It leads to a natural justification of chiral perturbation theory that works well in the structure of low-lying baryons as well as of nuclei. Being at a seminal stage, however, not much is known of the property of holographic-QCD skyrmions in medium, so we can say nothing of what that new skyrmion picture implies for dense baryonic matter. What we will do is to consider the skyrmion description where hidden local symmetry will be truncated so that only the lowest members of the tower ρ , ω and ϕ figure as relevant vector degrees of freedom (HLS $_{K=1}$ for short to be distinguished from HLS $_{K=\infty}$). Such an object is not well understood even at the level of one baryon, so our treatment of many-baryon systems will be qualitative at best.

The hidden local symmetry HLS $_{K=1}$ developed by Harada and Yamawaki is described in detail in Chapter 5. One can think of HLS $_{K=1}$ as a truncated model of HLS $_{K=\infty}$ in which all the members of the tower *except for the lowest* are integrated out. Quantum treatment of HLS $_{K=1}$ theory at one-loop level predicts what constitutes the principal mechanism of how the mass can disappear in heat bath or in dense medium. This is signalled at the chiral restoration phase transition – driven by temperature and/or density – by what is called “vector manifestation” (VM for short). At the VM fixed point, the vector meson mass, generically denoted by m_V , vanishes (when the quark masses are ignored corresponding to the “chiral limit”) proportionally to the quark condensate $\langle \bar{q}q \rangle$ which plays the role of a chiral order parameter. Although not established rigorously, the effect of the infinite tower remnant in HLS $_{K=1}$ can be interpreted in terms of the parameter a of the HLS $_{K=1}$ Lagrangian. The parameter a takes the value of 2 in free space but will approach 1 at the VM fixed point.

How the mass changes as the vacuum changes under extreme conditions of high temperature and high density is treated in Chapter 7. Here the VM fixed point of HLS $_{K=1}$ controls what can happen in relativistic heavy ion collisions where temperature can be hundreds of MeV. As for dense matter, the HLS $_{K=1}$ contains no explicit fermions, hence no baryons. They could in principle be generated as described in

Chapter 6 as skyrmions but in the absence of workable techniques, quasiquarks (or constituent quarks) are introduced in a hidden local symmetric way to $\text{HLS}_{K=1}$ Lagrangian. It is found that the VM fixed point is also reached when chiral symmetry is restored by density. In all cases, the vanishing of the mass m_V and a approaching 1 play an important role in the change of hadron properties in medium. How this intricate chiral symmetry manifestation, which appears different from the standard scenario exemplified by sigma models, affects nuclear matter at the normal nuclear matter density at zero temperature is discussed in Chapter 8. Here the scaling relation proposed a decade and half ago – called “Brown-Rho scaling” – makes a contact with the prediction of $\text{HLS}_{K=1}$ theory with the VM fixed point.

Strangeness at high density, in particular kaon condensation and possibly dense kaon bound nuclei, are treated from different vantage points in Chapter 9. A particular importance will be given to the approach that starts from the VM fixed point of hidden local symmetry theory, that makes a remarkably simple prediction that kaons will condense at a density ~ 3 times the nuclear matter density. This prediction will be the basis of the discussions on compact stars in Chapter 11. Other forms of dense matter that could figure at higher densities are discussed in Chapter 10. They are treated in the light of the “continuity” adopted in this volume. As such, it is not meant to be general. Nor can it directly confront the physics of compact stars as the intermediate steps are missing. In fact, there are generically two obstacles in confronting the physics of compact stars in terms of what’s discussed in this chapter. On the one hand, there is a vast variety of theoretical models in the literature that differ in details but cannot be precise enough to make falsifiable predictions. $\text{HLS}_{K=1}$, having been largely unexplored up to date in the density regime relevant for compact stars, is no exception in this regard. On the other hand, while there are a wealth of observations on properties of compact stars, they are not specific enough to confirm or falsify specific theoretical models. In other words, it is difficult to pin down what it is that a certain observation is telling us. Thus the discussion in Chapter 10, which addresses certain aspects of the most interesting density regime where perturbative QCD can make statements, cannot – and does not – directly address the properties of compact stars *per se*.

Some specific aspects of compact stars that bear on the equation of state of dense matter make up the content of Chapter 11. The focus will be on the neutron star mass which is the most important astrophysical observable that can be directly compared with theoretical results. In this chapter, an approach that departs drastically from conventional approaches found in the literature will be adopted. It will consist of assuming that the first phase transition that occurs beyond the nuclear matter density is kaon condensation and it determines the fate of compact stars, *i.e.* as to whether a star becomes a neutron star or collapses to a black hole. This would imply that other states at high densities discussed in Chapter 10 would be rendered irrelevant to compact stars.

Due to the limitation of the scope, we won't be able to adequately confront such related astrophysical issues as the formation and evolution processes of neutron stars, prospects for gravitational waves and gamma-ray bursts involving compact stars, neutron star cooling, and the neutrino signals in the process of neutron star formation. Even though these observations may not be able to provide directly details of the inside structure of neutron stars, they will certainly supply very useful information on masses and populations of compact objects.

The conclusion we can draw at this point *vis-à-vis* with compact stars is that a lot more work is needed and a great deal of discovery, both theoretical and experimental, are in store. The penultimate goal of probing dense hadronic matter via compact stars remains a challenge for the future.

In exposing the principal structure of this volume, we will take for granted some of the basic premises that are treated in CND-I. For completeness, we make a mini-summary of the basic elements that we will have in mind in this volume. They will be stated without explanation.

Confinement

We will take it for granted that at zero temperature and zero density, quarks and gluons are confined and that at some temperature $T_{c1} > 0$ and some density $n_{c1} > n_0$ where $n_0 \simeq 0.16 \text{ fm}^{-3}$ is the normal nuclear matter density, quarks and gluons are deconfined. How this happens which we take as unknown will not be addressed.

Chiral symmetry

We will be mainly dealing with the up (u) and down (d) quarks but will consider the strange (s) quark as well. We will be mostly considering the "chiral limit" where $m_u \approx m_d \approx m_s \approx 0$ although in nature, the u- and d-quark masses are a few MeV and the s-quark mass ~ 100 MeV. Unless otherwise stated, these will be referred to as "light quarks." The QCD Lagrangian is invariant under chiral transformations. This is the $SU(N_f)_L \times SU(N_f)_R$ chiral symmetry.¹ Thus both the left (L) and right (R) currents or alternatively the vector and axial vector currents are conserved. The conservation of the vector current is not affected if the quark masses are turned on provided the vector ($V = L + R$) symmetry is unbroken but the axial current will have a term proportional to the quark mass if the quark mass is not zero.

This symmetry of Lagrangian can, however, be broken by the vacuum. It turns out that the vector symmetry remains unbroken by the vacuum (Vafa-Witten theorem, see below) but the axial symmetry is broken, as we know from Goldstone theorem. It is now understood – and supported by nature – that the symmetry associated with the light quarks in the zero-mass limit is spontaneously broken below some temperature T_{c2} and some density n_{c2} . The present understanding is $T_{c2} \simeq 170 - 190$ MeV and $n_{c2} \simeq (3 - 10)n_0$. The latter is more or less unknown.

¹There are in addition an unbroken $U(1)_V$ symmetry corresponding to baryon charge and anomalous $U(1)_A$ which will be mentioned below.

Whether the critical temperature and density (T_{c_2}, n_{c_2}) for chiral restoration are the same as or different from the deconfinement points (T_{c_1}, n_{c_1}) is not known. We will encounter the situation in which deconfinement may take place before chiral symmetry is restored, *i.e.* $T_{c_1} < T_{c_2}$. There are no known models that indicate the converse takes place. Unless otherwise stated, we shall take them coincident. Lattice results are not at odds with this choice. It will be taken for granted that the $SU(N_f) \times SU(N_f)$ chiral symmetry is spontaneously broken to the diagonal subgroup $SU(N_f)_{V=L+R}$. Thus we will have $(N_f^2 - 1)$ Goldstone (or with quark masses, pseudo-Goldstone) bosons, π and a massive η' .

Anomalies

Certain symmetries of QCD present in the “bare” Lagrangian are broken by quantum effects intimately tied to the complex vacuum structure. Such a phenomenon breaks the current conservation unlike the spontaneous breaking which leaves the relevant currents conserved. This is what is called “quantum anomaly” or “anomaly” in short. Among a variety of anomalies associated with QCD proper and effective theories thereof, what concerns us are the scale invariance and the axial $U(1)$ symmetry in the limit of zero mass.

- *Trace anomaly:*

The QCD Lagrangian with no quark masses is scale invariant. So the energy-momentum tensor $\theta_{\mu\nu}$ is traceless at the classical level. However at the quantum level due to an anomaly, scale invariance is broken so the trace of the energy-momentum tensor has a non-zero value² $\theta_{\mu}^{\mu} = -\frac{\beta(g)}{2g} \text{Tr} G_{\mu\nu}^2$. The non-vanishing trace of $\theta_{\mu\nu}$ means that there is a scale in the theory, so the scale invariance present in the Lagrangian is broken. If one defines a dilation current D_{μ} , the scale anomaly means that the dilatation current is not conserved. The divergence of the dilatation current is equal to the trace of $\theta_{\mu\nu}$.

We will later observe, based on lattice data, that the vacuum expectation value of $G_{\mu\nu}^2$, *i.e.* $\langle 0|G^2|0\rangle$, could be decomposed into two components “soft” and “hard,” $\langle 0|G^2|0\rangle = \langle 0|G^2|0\rangle_{soft} + \langle 0|G^2|0\rangle_{hard}$. The “soft” component will be interpreted later as due to spontaneous breaking of scale invariance and the “hard” component to explicit breaking of scale invariance. It is understood, though not rigorously proven, that spontaneous breaking of the scale invariance is possible only if there is explicit breaking. This feature will play an important role in some parts of this volume.

- $U(1)_A$ anomaly

While the $SU(N_f)$ axial currents are conserved, spontaneously broken in low temperature and low density regime, and assumed to be restored at T_{c_2}

²With non-zero quark mass there is an additional term on the RHS of the form $+\sum_f m_f(1 + \gamma_f)\bar{q}_f q_f$ with γ_f the anomalous dimension for the quark field q_f .

and n_{c2} , the $U(1)$ axial current is not conserved due to a triangle anomaly. The consequence is that the would-be Goldstone boson η' is not massless even in the chiral limit. The anomaly is believed to remain in the chirally restored phase. In this volume, this anomaly will not figure importantly.

Large N_c

QCD in the large N_c limit provides useful information on nonperturbative aspects of the strong interactions. In nature, N_c is 3 but in some sense which cannot be given a precise meaning, $N_c = 3$ can be considered much bigger than the number of flavors $N_f = 2, 3$. In terms of the power of N_c factors, the following counting can be done:

- *Gluons*

The counting in the gluon sector is done in terms of the 't Hooft constant $\lambda = g^2 N_c$. This is taken to be $\mathcal{O}(1)$.³ Thus $g^2 \sim \mathcal{O}(1/N_c^2)$. Planar diagrams dominate, with non-planar graphs being suppressed by $1/N_c$ or higher.

- *Mesons*

The leading contribution to correlation functions is the sum of planar graphs which is $\mathcal{O}(N_c)$. Meson-meson couplings are down by powers of $1/N_c$. For instance, 3-meson interactions scale as $1/\sqrt{N_c}$, 4-meson interactions $1/N_c$ etc. The mass of the $\bar{q}q$ -type mesons is $\mathcal{O}(N_c^0)$. They appear as sharp resonances. Exotics are suppressed by $\mathcal{O}(1/N_c)$ or higher.

- *Baryons*

Baryon mass scales as $\mathcal{O}(N_c)$. This follows from the fact that N_c quarks are needed to make a “color” singlet. This plays an important role in what follows. Roughly

$$m_B \sim N_c \sim 1/g^2 \tag{1.1}$$

where g is the color gauge coupling. The mass being inversely proportional to g^2 indicates that the baryon can be considered as a soliton in a weak coupling theory. Indeed the skyrmion description of baryons is intimately tied to this scaling as we will see later. Simple counting shows that $g_A \sim N_c$, $f_\pi \sim \sqrt{N_c}$ and $g_{\pi NN} \sim N_c^{3/2}$. There are a variety of relations between the skyrmion model and non-relativistic quark models in the large N_c limit which are listed in CND-I, which we will use without entering into details.

³In holographic QCD discussed in several chapters that follow, λ will be taken to be large in the sense that an expansion is made in $1/\lambda$.

Vafa-Witten theorem

Modulo some plausible assumptions, vector symmetries in vector-like theories like QCD are argued to be free from spontaneous symmetry breaking. The examples relevant to us are baryon number, isospin *etc.* One of the assumptions is that Lorentz invariance holds. Since the systems we will be considering in this volume will break explicitly Lorentz invariance (by temperature and/or density), this Vafa-Witten theorem does not directly apply, although it will be a useful guide in writing down the Lorentz-invariant part of effective theories.

Anomaly matching

In order for an effective theory to be consistent with a fundamental theory, the anomaly in EFT should be identical to the anomaly in the fundamental theory. This is 't Hooft anomaly matching condition. In QCD, this means that in the hadronic world where hadrons are composites of quarks and gluons, there must be a massless, colorless object which couples to the axial current and photon and reproduces exactly the triangle anomaly. The Goldstone pion is such an object. This condition will prove to be quite relevant in the chiral structure of the nucleon treated in Chapter 8.