Coulomb Gases and Ginzburg-Landau Vortices

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March 28, 2014
Acknowledgments
These are the lecture notes of the “Nachdiplomvorlesung” course that I taught in the Spring of 2013 at the ETH Zürich, at the invitation of the Forschungsinstitut für Mathematik. I would like to express my deep gratitude to Tristan Rivièire, director of the FIM, for this invitation, and to him, Thomas Kappeler, Michael Struwe and Francesca da Lio, for making my stay particularly enjoyable and stimulating.

I am grateful to those who followed the course for their constructive questions and comments. A particular acknowledgment goes to Thomas Leblé for taking notes and typing the first draft of the manuscript, as well as for making the figures.

The material covered here is based on works with Etienne Sandier and with Nicolas Rougerie, I am very grateful to them for the fruitful collaborations we have had. Finally, this manuscript also benefited from comments by Ofer Zeitouni to whom I extend thanks here.
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1 Introduction

These lecture notes are devoted to the mathematical study of two physical phenomena that have close mathematical connections: vortices in the Ginzburg-Landau model of superconductivity on the one hand, and classical Coulomb gases on the other hand. A large part of the results we shall present originates in joint work with Etienne Sandier (for Ginzburg-Landau and two dimensional Coulomb gases) and in joint work with Nicolas Rougerie (for higher dimensional Coulomb gases). In order to simplify the presentation, we have chosen to present the material in reverse chronological order, starting with the more recent results on Coulomb gases which are simpler to present, and finishing with the more complex study of vortices in the Ginzburg-Landau model. But first, in this introductory chapter, we start by briefly presenting the two topics, the context and the connection between them.

1.1 From the Ginzburg-Landau model to the 2D Coulomb gas

1.1.1 Superconductivity and the Ginzburg-Landau model

The Ginzburg-Landau model is a very famous physics model for superconductivity. Superconductors are certain metallic alloys, which, when cooled down below a very low critical temperature, lose their resistivity and let current flow without loss of energy. This is explained at the microscopic level by the creation of superconducting electron pairs called Cooper pairs (Bardeen-Cooper-Schrieffer or BCS theory), and superconductivity is a macroscopic manifestation of this quantum phenomenon. The Ginzburg-Landau theory, introduced on phenomenological grounds by Ginzburg and Landau in the 1950’s, some forty years after superconductivity had first been discovered by Kammerling Ohnes in 1911, has proven amazingly effective in describing the experimental results and predicting the behavior of superconductors. It is only very recently that the Ginzburg-Landau theory [GL] has been rigorously (mathematically) derived from the microscopic theory of Bardeen-Cooper-Schrieffer [BCS], also dating from the 50’s, by Frank, Hainzl, Seiringer and Solovej [FHSS].

These superconducting alloys exhibit a particular behavior in the presence of a magnetic field: the superconductor levitates above the magnet. This is explained by the Meissner effect: the superconductor expells the magnetic field. This only happens when the external field $h_{ex}$ is not too large. There are three critical fields $H_{c1}, H_{c2}, H_{c3}$ for which phase transitions occur. Below the first critical field $H_{c1}$, the material is everywhere superconducting. At $H_{c1}$, one first observes local de-
flicts of superconductivity, called vortices, around which a superconducting loop of current circulates. As $h_{\text{ex}}$ increases, so does the number of vortices, so that they become densely packed in the sample. The vortices repel each other, while the magnetic field confines them inside the sample, and the result of the competition between these two effects is that they arrange themselves in a particular triangular lattice pattern. It was predicted by Abrikosov [Abri], and later observed experimentally, that there should be periodic arrays of vortices appearing in superconductors, and this was later observed experimentally (Abrikosov and Ginzburg earned the 2003 Nobel Prize for their discoveries on superconductivity), cf. Fig. 1.1 (for more pictures, see www.fys.uio.no/super/vortex/).

These triangular lattices (originally Abrikosov predicted a square lattice but he had made a small mistake) then became called Abrikosov lattices. A part of our study, detailed in this course, is aimed towards understanding why this particular lattice appears.

The second and third critical fields correspond respectively to the loss of superconductivity in the sample bulk and to the complete loss of superconductivity. These two transitions are not the focus of our study, and for more mathematical details on them, we refer to the monograph by Fournais and Helffer [FH1]. For a physics presentation of superconductivity and the Ginzburg-Landau model we refer to the standard texts [SST, DeG, Ti], for a mathematical presentation one can see [SS4, FH1] and references therein.

In non-dimensionalized form and in a simply connected domain $\Omega$ of the plane, the model proposed by Ginzburg-Landau can be written as the functional

$$G_\varepsilon(u,A) = \frac{1}{2} \int_\Omega |(\nabla - iA)u|^2 + |\text{curl} A - h_{\text{ex}}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2}. \quad (1.1)$$

This may correspond to the idealized situation of an infinite vertical cylindrical
1.1 From the Ginzburg-Landau model to the 2D Coulomb gas

A sample of cross-section Ω and a vertical external field of intensity \( h_{\text{ex}} \), or to a thin film.

Here

- \( u : \Omega \to \mathbb{C} \), usually denoted by \( \psi \) in the physics literature, is called the order parameter. Its modulus (the density of Cooper pairs of superconducting electrons in the BCS theory) indicates the local state of the material: where \( |u| \approx 1 \), the material is in the superconducting phase, where \( |u| \approx 0 \) in the normal phase. The vortices correspond to isolated zeroes of \( u \), and since \( u \) is complex-valued each zero carries an integer topological degree, like a "topological charge."

- \( A : \Omega \to \mathbb{R}^2 \) is the vector potential of the magnetic field \( h = \text{curl} \ A \) (defined by \( \text{curl} \ A := \partial_1 A_2 - \partial_2 A_1 \)), which is thus a real-valued function.

- The parameter \( h_{\text{ex}} > 0 \) is the intensity of the applied (or external) magnetic field.

- The parameter \( \varepsilon > 0 \) is a material constant, corresponding to the ratio between characteristic length scales of the material (the coherence length over the penetration depth). We will be interested in the asymptotic regime \( \varepsilon \to 0 \). The functional is generally expressed in the physics literature in terms of the inverse of the constant \( \varepsilon \), denoted \( \kappa \), and called the Ginzburg-Landau parameter. Materials with high-\( \kappa \) (the case we are interested in) are sometimes called "extreme type-II superconductors," and the limit \( \kappa \to \infty \) is often called the London limit.

When considering the problem of minimizing the functional \( G_\varepsilon \), a heuristic examination leads to observing that:

- The term \((1 - |u|^2)^2\) favors \( u \) close to 1, hence \( u \) should not vanish too often, especially as \( \varepsilon \to 0 \). A dimensional analysis in fact shows that the regions where \( |u| \) is small have lengthscale \( \varepsilon \).

- The quantity \(|\text{curl} \ A - h_{\text{ex}}|^2\) is smaller when \( \text{curl} \ A = h \approx h_{\text{ex}} \), that is, when the magnetic field penetrates the material so that the induced magnetic field equals the external magnetic field.

Minimizers and critical points of the Ginzburg-Landau functional without boundary constraints solve the associated set of Euler-Lagrange equations, called the Ginzburg-Landau equations:

\[
\begin{aligned}
\text{(GL)} \quad &\left\{ \begin{array}{l}
-\nabla_A^2 u + \frac{1}{\varepsilon^2} u (1 - |u|^2) \quad \text{in } \Omega \\
-\nabla^\perp h = \langle iu, \nabla_A u \rangle \quad \text{in } \Omega
\end{array} \right.
\end{aligned}
\]

where \( \nabla_A := \nabla - iA \), \( \langle \cdot, \cdot \rangle \) denotes the scalar product in \( \mathbb{C} \) as identified with \( \mathbb{R}^2 \), \( \nabla^\perp = (-\partial_2, \partial_1) \) and again \( h = \text{curl} \ A \); with natural boundary conditions

\[
\begin{aligned}
\nabla_A u \cdot \nu &= 0 \quad \text{on } \partial \Omega \\
h &= h_{\text{ex}} \quad \text{on } \partial \Omega.
\end{aligned}
\]
1.1.2 Reduction to a Coulomb interaction

More details on the analysis of the Ginzburg-Landau model will be given in Chapter 7, which will be devoted to it, but for now let us try to explain the Coulombic flavor of the phenomenon.

In the regime with vortices (for $H_{c_1} \leq h_{ex} \ll H_{c_2}$), formal computations that will be better detailed in Chapter 7 show that in the asymptotic regime $\varepsilon \to 0$, the functional $G_\varepsilon(u, A)$ behaves as if it were:

$$G_\varepsilon(u, A) \approx \frac{1}{2} \int_\Omega |\nabla h|^2 + |h - h_{ex}|^2,$$

where $h = \text{curl} A$ (1.2)

with what is known in the physics literature as the London equation:

$$\begin{cases}
-\Delta h + h \approx 2\pi \sum_i d_i \delta^{(c)}_{a_i} & \text{in } \Omega \\
h = h_{ex} & \text{on } \partial \Omega,
\end{cases}$$

(1.3)

where the $a_i$'s are the centers of the vortices of $u$ and the coefficients $d_i \in \mathbb{Z}$ their (topological) degrees. One should think of $\delta^{(c)}_{a_i}$ as being formally a Dirac mass at $a_i$, smoothed out at the scale $\varepsilon$, or some approximation of it. A large part of our analysis in [SS4,SS7] is devoted to giving rigorous statements and proofs of these heuristics.

Inserting the London equation (1.3) into the approximation (1.2) leads to the following electrostatic analogy:

$$G_\varepsilon(u, A) \approx \frac{1}{2} \int_{\Omega \times \Omega} G_\Omega(x, y) \left( 2\pi \sum_i d_i \delta^{(c)}_{a_i} - h_{ex} \right)(x) dx \left( 2\pi \sum_i d_i \delta^{(c)}_{a_i} - h_{ex} \right)(y) dy$$

(1.4)

where $G_\Omega$ is a Green kernel (or more accurately, Yukawa or screened Green kernel), solution to

$$\begin{cases}
-\Delta G_\Omega + G_\Omega = \delta_y & \text{in } \Omega \\
G_\Omega = 0 & \text{on } \partial \Omega.
\end{cases}$$

(1.5)

This kernel is logarithmic to leading order: we may write

$$G_\Omega(x, y) = -\frac{1}{2\pi} \log |x - y| + R_\Omega(x, y)$$

(1.6)

where $R_\Omega$ is a nonsingular function of $(x, y)$. Approximating $G_\Omega$ by $-\frac{1}{2\pi} \log$ gives that the leading terms in (1.4) are

$$G_\varepsilon(u, A) \approx -\pi \sum_{i,j} d_i d_j \log |a_i - a_j|$$

(1.7)

which is a sum of pairwise logarithmic or Coulombic interactions, weighted by the degrees $d_i$. Two such topological charges repel each other when they have the same sign, and attract each other if they have different signs. Rigorously, this is of course wrong, because we have replaced the smoothed out Diracs by true Dirac masses,
leading to infinite contributions when \( i = j \) in (1.7). One needs to analyze more carefully the effects of the smearing out, and to remove the infinite self-interaction of each “charge” at \( a_i \) in (1.7). One also needs to retain the interaction of these charges with the “background charge” \(-h_{ex} dx\) appearing in (1.4). This is what leads to the analogy with the Coulomb gas that we will define and describe just below.

When looking for a model that retains these features: Coulombic interactions of points, combined with the confinement by a background charge, the simplest is to consider a discrete model with all charges equal to 1, and consider the Hamiltonian of a Coulomb gas with confining potential in dimension 2:

\[
H_n(x_1, \ldots, x_n) = -\sum_{i \neq j} \log |x_i - x_j| + n \sum_{i=1}^n V(x_i),
\]

where \( x_i \in \mathbb{R}^2 \), \( V \) is the confining potential (smooth, growing faster than \( \log |x| \) at infinity), and the number of points \( n \) tends to infinity.

It turns out that this much simpler model (compared to \( G_\varepsilon \)) does retain many of the essential features of the vortex interaction, and is also of independent interest for physics and mathematics, as we will see. The study of (1.8) and its higher-dimensional analogues will occupy the largest part of these notes. We will then see how to use the perspective and knowledge gained on this to analyze the Ginzburg-Landau model (again, this is the reverse of the literature chronology, since we first studied the Ginzburg-Landau model and then adapted our analysis to the Coulomb gas situation!).

1.2 The classical Coulomb gas

1.2.1 The general setting

The Hamiltonian given by (1.8) corresponds to the energy of a gas of charged particles in \( \mathbb{R}^2 \) interacting via the Coulomb kernel in two dimension. To be more precise, \(-\log |x - y|\) is a multiple of the Coulomb kernel (or the fundamental solution of the Laplacian in the plane) in dimension 2. The counterpart in higher dimension corresponds to the \( d \)-dimensional Coulomb kernel, which is a multiple of \(|x|^{2-d}\) for \( d \geq 3 \). The Hamiltonian of a classical Coulomb gas in any dimension \( d \geq 2 \) is thus given by

\[
H_n(x_1, \ldots, x_n) = \sum_{i \neq j} g(x_i - x_j) + n \sum_{i=1}^n V(x_i),
\]

where

\[
g(x) = \begin{cases} 
-\log |x| & \text{for } d = 2 \\
\frac{1}{|x|^{d-2}} & \text{for } d \geq 3.
\end{cases}
\]
The statistical mechanics of a Coulomb gas is described by the corresponding
Gibbs measure:
\[
\frac{dP_{n,\beta}}{Z_{n,\beta}}(x_1, \ldots, x_n) := \frac{1}{Z_{n,\beta}} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n \tag{1.11}
\]
where \(\beta > 0\) is the inverse temperature and \(Z_{n,\beta}\) is a normalization constant, the
partition function, defined by
\[
Z_{n,\beta} = \int_{(\mathbb{R}^d)^n} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n.
\]

The probability measure \(P_{n,\beta}\) gives the probability of finding the particles at
\((x_1, \ldots, x_n)\) at (inverse) temperature \(\beta\). The object of statistical mechanics is then
to analyze possible transitions in the types of states that can be effectively
observed (i.e. those that have probability 1 or almost 1), according to the value
of the inverse temperature \(\beta\) (e.g. transitions from ordered to disordered states at
critical temperatures, such as liquid to solid phases etc). For general reference, we refer to standard statistical mechanics textbooks such as [Huan], and for the more
specific Coulomb gas systems, to the book by Forrester [Fo].

This model is one of the most basic statistical mechanics models not confined
to a lattice, and it is considered difficult because of the long-range nature of the
electrostatic interaction. Moreover, it can play the role of a toy model of the
structure of matter, even if it is a purely classical - and not quantum - model.
Studies in this direction include [SM, LieOx, JLM, PenSm].

The macroscopic distribution of the points as their number \(n\) goes to infinity
is well understood and , relatively simple to derive, this will be the object of
Chapter 2. On the other hand, their microscopic distribution, more precisely the
one seen at the scale \(n^{-1/d}\), is less understood, and will be the main object of these
lectures.

Let us now see some more specific motivations for studying the classical Coulomb
gas, many of them being specific to dimension \(d = 2\).

### 1.2.2 Two-dimensional Coulomb gas

This is the setting that is the closest to the Ginzburg-Landau setting, as we
discussed above. In this setting, the microscopic distribution of the points in the
plane is expected in some regimes to follow Abrikosov lattice triangular patterns.

**Vortices in superfluids and superconductors** A first motivation for studying
the two-dimensional Coulomb gas is the analysis of vortices in the Ginzburg-
Landau model of superconductivity, but also more generally of vortex systems in
classical fluids [CLMP], in quantum fluids such as in superfluids or Bose-Einstein
condensates [CPRY], and in fractional quantum Hall physics [Gir, RSY1, RSY2].
All these systems share a lot of mathematics in common, and it is also of interest to
understand their statistical physics (critical temperatures and phase transitions).
1.2 The classical Coulomb gas

Fekete sets  This motivation no longer comes from physics but rather from a very different area of mathematics: interpolation theory. Fekete points are defined to be points that maximize the quantity

$$\prod_{1 \leq i < j \leq n} |x_i - x_j|,$$

among all families of $n$ points defined on a certain subset of $\mathbb{R}^d$, or a manifold (or any metric space, replacing modulus by the distances). The Fekete points have the property of minimizing the error when interpolating a function by its value at points, see [SaTo] for reference, or [SK] for more details on the motivation. A whole literature is also devoted to understanding Fekete points on complex manifolds, possibly in higher dimension, see e.g. [Ber, BBN, LevOrC] and references therein.

Of course, in the setting of the Euclidean space, maximizing (1.12) is equivalent to minimizing the logarithmic interaction

$$-\sum_{1 \leq i \neq j \leq n} \log |x_i - x_j|$$

which takes us back to the setting of the two-dimensional Coulomb gas. Indeed, Fekete sets confined to a set $K \subset \mathbb{R}^d$ correspond to minimizers of $H_n$ with $V$ taken to be 0 in $K$ and $+\infty$ in $K^c$. Minimizers of (1.8) for general $V$’s are in fact called weighted Fekete sets, also defined as maximizers of

$$\prod_{1 \leq i < j \leq n} |x_i - x_j| \prod_{i=1}^n e^{-\frac{1}{2}V(x_i)}$$

where $V$ is the weight. For definitions and the connection to logarithmic potential theory, see again [SaTo] and references therein. Weighted Fekete sets are also naturally related to the theory of weighted orthogonal polynomials (cf. the surveys [Sim, Ko] or again [SaTo]).

The correspondence can also be made via a mapping, e.g. the important question of finding the Fekete points on the (2-)sphere is equivalent, by stereographical projection, to studying the weighted Fekete sets on $\mathbb{R}^2$ with weight $V(x) = \frac{1}{2} \log(1 + |x|^2)$, for details see [Ha].

Random matrix theory  Random matrix theory (RMT) is a relatively old theory, pioneered by statisticians and physicists such as Wishart, Wigner and Dyson, and originally motivated by the understanding of the spectrum of heavy atoms, see [Ma]. For more recent mathematical reference see [ACZ, D, Fo]. An important model of random matrices is the so-called Ginibre ensemble [Gin]: the law is that of an $n \times n$ complex matrix whose coefficients are i.i.d. complex normal random variables. The main question asked by RMT is: what is the law of the spectrum of a large random matrix? In the case of the Ginibre ensemble, the law is known exactly: upon rescaling the (complex) eigenvalues $x_1, \ldots, x_n$ by a factor
\( \frac{1}{\sqrt{n}} \), it is given by the following density

\[
d\mathcal{P}_n(x_1, \ldots, x_n) = \frac{1}{Z_n} e^{-H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n
\]

(1.13)

with

\[
H_n(x_1, \ldots, x_n) = - \sum_{i \neq j} \log |x_i - x_j| + n \sum_{i=1}^{n} |x_i|^2
\]

(1.14)

and \( Z_n \) a normalization constant. We recognize in (1.14) the 2D Coulomb gas Hamiltonian with potential \( V(x) = |x|^2 \), and the law \( d\mathcal{P}_n \) is the Gibbs measure (1.11) at inverse temperature \( \beta = 2 \). This analogy between random matrices and the statistical mechanics of Coulomb gases was first noticed by Wigner [Wi] and Dyson [Dy], see [Fo] for more on this link. Writing the law in the form (1.13) immediately displays the phenomenon of repulsion of eigenvalues: eigenvalues in the complex plane interact like Coulomb particles, i.e. they do not “like” to be too close and repel each other logarithmically.

At this specific temperature \( \beta = 2 \), the law of the spectrum acquires a special algebraic feature: it becomes a determinantal process, part of a wider class of processes (see [HKPV,Bor]) for which the correlation functions are explicitly given by certain determinants. This allows for many explicit algebraic computations. However, many relevant quantities that can be computed explicitly for \( \beta = 2 \) are not exactly known for the \( \beta \neq 2 \) case, even in the case of the potential \( V(x) = |x|^2 \).

In this course, in contrast, we will work for any \( \beta \), and with a wide class of potentials.

### 1.2.3 The one-dimensional Coulomb gas and the log gas

We have not mentioned yet the one-dimensional Coulomb gas, which corresponds to (1.9) with the Coulomb kernel (up to a constant) \( g(x) = |x| \). The reason is that we will not be interested in it, because it has already been well-understood [Le1, Le2, Ku, BraLie, AlMu]. It can be “solved” almost explicitly and crystallization at zero temperature is established.

We are interested however in another one-dimensional model (i.e. with points \( x_i \in \mathbb{R} \)), where the two-dimensional logarithmic interaction \( g(x) = -\log |x| \) is used in (1.9). This is usually called a log gas, and its motivation also comes from Random Matrix Theory (see [Fo]): one-dimensional counterparts to the Ginibre ensemble are the Gaussian Unitary Ensemble (GUE) and the Gaussian Orthogonal Ensemble (GOE), which are symmetric analogues of it. The law of the GUE (resp. the GOE) is that of a \( n \times n \) matrix whose coefficients are complex (resp. real) normal random variables, independent up to a Hermitian (resp. symmetry) condition. Because of the Hermitian or symmetric nature of the matrix, its eigenvalues lie on the real line (hence the one dimensionality of the model), but they still repel each other logarithmically: again, the law of the spectrum (the distribution of
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eigenvalues) can be given explicitly by the following density on $\mathbb{R}$
\[
d\Pi_n(x_1, \ldots, x_n) = \frac{1}{Z_n} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n,
\]
where $H_n$ is still defined as
\[
H_n(x_1, \ldots, x_n) = -\sum_{i \neq j} \log |x_i - x_j| + n \sum_{i=1}^n |x_i|^2, \quad x_i \in \mathbb{R}
\]
with $\beta = 1$ for the GOE and $\beta = 2$ for GUE. This is thus a particular case of a \textit{log gas}, at specific temperature $\beta = 1$ or $2$ and with quadratic potential, and the phenomenon of \textit{repulsion of eigenlevels} is visible in the same way as for the Ginibre ensemble. Again, in these cases of the GOE and GUE, a lot about \ref{1.15} can be understood and computed explicitly thanks to the underlying random matrix structure and its determinantal nature. In fact the global and local statistics of eigenvalues are completely understood.

Considering the coincidence between a statistical mechanics model and the law of the spectrum of a random matrix model for several values of the inverse temperature, it is also natural to ask whether such a correspondence exists for any value of $\beta$. The answer is yes for $\beta = 4$, it corresponds to the Gaussian Symplectic Ensemble (GSE) of Hermitian matrices with quaternionic coefficients, and for any $\beta$ a somehow complicated model of tridiagonal matrices can be associated to the Gibbs measure of the one-dimensional log gas at inverse temperature $\beta$, see \cite{DE}. This and other methods allow again to compute a lot explicitly, and to derive that the microscopic laws of the eigenvalues are those of a so called \textit{sine-$\beta$ process} \cite{VV}.

Generally speaking, much is known for log gases in one dimension, for any value of $\beta$ and a wide class of potentials $V$ (see \cite{BEY1, BEY2}). In contrast, the analogue is true in dimension $2$ only for $\beta = 2$ \cite{Gin}. Thus the topic of log/Coulomb gases does not seem reducible to just a subset of Random Matrix Theory. This is of course even more true in dimension $3$ and higher, which we will also treat, and where we leave the realm of RMT.
2 The leading order behavior of the Coulomb gas

In this chapter, we study the leading order or “mean field" behavior of the Coulomb gas Hamiltonian. The results are quite standard and borrowed or adapted from the literature. However, we try to give here a self-contained and general treatment, since results in the literature are a bit scattered between the potential theory literature, the probability and statistical mechanics literature and the PDE literature, and not all situations seem to be systematically covered.

Let us recall the setting : for \((x_1, \ldots, x_n) \in (\mathbb{R}^d)^n\), we define the Hamiltonian

\[
H_n(x_1, \ldots, x_n) = \sum_{i \neq j} g(x_i - x_j) + n \sum_{i=1}^n V(x_i)
\]  

(2.1)

where \(g\) is a multiple of the Coulomb kernel in dimension \(d \geq 2\) (the fundamental solution of the Laplacian) and is \( -\log\) in dimension 1 (this choice is made to treat the case of log gases mentioned in the previous chapter):

\[
g(x) = \begin{cases} 
-\log |x| & \text{for } d = 1, 2 \\
\frac{1}{|x|^{d-2}} & \text{for } d \geq 3.
\end{cases}
\]  

(2.2)

We need to keep track of the proportionality factor between \(g\) and the true Coulomb kernel, and note that, for \(d \geq 2\), we have

\[-\Delta g = c_d \delta_0
\]  

(2.3)

where \(\delta_0\) is the Dirac mass at \(0 \in \mathbb{R}^d\), and where the constant \(c_d\) is given by :

\[c_2 = 2\pi \quad \text{and } c_d = (d - 2)|\mathbb{S}^{d-1}| \text{ for } d \geq 3,
\]  

(2.4)

see e.g. [LiLo, Chap. 6]. In dimension \(d = 1\), instead of (2.3), \(g\) solves the non-local equation

\[-\Delta^{1/2} g = c_1 \delta_0 \quad c_1 = \pi,
\]

where \(\Delta^s\) is the fractional Laplacian (in this situation one can check that \(\Delta^{1/2}\) coincides with the Dirichlet-to-Neumann operator in the upper half-plane, see e.g. [CaRS]). We will sometimes abuse notation by considering also \(g\) as a function on \(\mathbb{R}\), by writing \(g(r) = \begin{cases} 
-\log r & \text{for } d = 1, 2 \\
\frac{r^{2-d}}{\pi^{d-2}} & \text{for } d \geq 3.
\end{cases}
\)  

The function \(V\) is called the (confining) potential, precise assumptions on \(V\) will be made later. The Hamiltonian \(H_n\) may be physically interpreted as follows : it is the sum of an interaction term \(\sum_{i \neq j} g(x_i - x_j)\) and a confining term \(n \sum_{i=1}^n V(x_i)\). The first term describes the pairwise interaction of charged particles of same sign,
interacting via (a multiple of) the Coulomb potential in dimension $d$. Since all
the charges have the same sign, their natural behavior is to repel each other, and
potentially escape to infinity. The potential $V$, however, is there to confine the
particles to a compact set of $\mathbb{R}^d$. Note that the sum of pairwise interactions is
expected to scale like the number of pairs of points, i.e. $n^2$, while the sum of the
potential terms is expected to scale like $n$ times the number of points, i.e. $n^2$
again. The factor $n$ in front of $V$ in (2.1) is there precisely so that this happens,
in such a way that the opposing effects of the repulsion and of the confinement
balance each other. This is called the “mean-field scaling”. It is the scaling in which
the force acting on each particle is given in terms of the average field generated
by the other particles. For general reference on mean-field theory, see statistical
mechanics textbooks such as [Huan].

The beginning of this chapter is devoted to the analysis of $H_n$ only, via Γ-
convergence, leading to the mean-field description of its minimizers. Later, in
Section 2.6 we apply these results to the statistical mechanics model associated to
the Hamiltonian $H_n$, i.e. to characterizing the states with nonzero temperature.

### 2.1 Γ-convergence: general definition

The result we want to show about the leading order behavior of $H_n$ can be
formalized in terms of the notion of Γ-convergence, in the sense of De Giorgi
(see [Braides1], [Braides2] for an introduction, or [DalM] for an advanced refer-
ence). It is a notion of convergence for functions (or functionals) which ensures
that minimizers tend to minimizers. This notion is popular in the community
of calculus of variations and very used in the analysis of sharp-interface and frac-
ture models, dimension reduction for variational problems, homogeneization... (see
again [Braides1] for a review) and even more recently in the study of some evolution
problems [Braides3]. Using this formalism here is convenient but not essential.

Let us first give the basic definitions.

**Definition 2.1** (Γ-convergence). We say that a sequence $\{F_n\}_n$ of functions on a
metric space $X$ Γ-converges to a function $F : X \to (-\infty, +\infty]$ if the following two
inequalities hold:

1. (Γ-lim inf) If $x_n \to x$ in $X$, then $\liminf_{n \to +\infty} F_n(x_n) \geq F(x)$.
2. (Γ-lim sup) For all $x$ in $X$, there is a sequence $\{x_n\}_n$ in $X$ such that $x_n \to
   x$ and $\limsup_{n \to +\infty} F_n(x_n) \leq F(x)$. Such a sequence is called a recovery
   sequence.

The second inequality is essentially saying that the first one is sharp, since
it implies that there is a particular sequence $x_n \to x$ for which the equality
$\lim_{n \to +\infty} F_n(x_n) = F(x)$ holds.
Remark 2.2. In practice a compactness assumption is generally needed and sometimes added in the definition, requiring that if \( \{ F_n(x_n) \} \) is bounded, then \( \{ x_n \} \) has a convergent subsequence. A similar compactness requirement also appears in the definition of a good rate function in large deviations theory (see Definition 2.24 below).

Remark 2.3. The first inequality is usually proven by functional analysis methods, without making any “ansatz” on the precise form of \( x_n \), whereas the second one is usually obtained by an explicit construction, during which one constructs “by hand” the recovery sequence such that that \( F_n(x_n) \) has asymptotically less energy than \( F(x) \). Note also that by a diagonal argument, one may often reduce to constructing a recovery sequence for a dense subset of \( x \)’s.

Remark 2.4. A \( \Gamma \)-limit is always lower semi-continuous (l.s.c.) as can be checked. (In particular, a function which is not l.s.c. is a bad candidate for being a \( \Gamma \)-limit.) Thus, a functional is not always its own \( \Gamma \)-limit : in general \( \Gamma \lim F = \bar{F} \) where \( \bar{F} \) is the l.s.c. envelope of \( F \).

Remark 2.5. The notion of \( \Gamma \)-convergence can be generalized to the situation where \( F_n \) and \( F \) are not defined on the same space. One may instead refer to a sense of convergence of \( x_n \) to \( x \) which is defined via the convergence of any specific function of \( x_n \) to \( x \), which may be a nonlinear function of \( x_n \), cf. [SS3, JSt] for instances of this.

We now state the most important property of \( \Gamma \)-convergence : \( \Gamma \)-convergence sends minimizers to minimizers.

Proposition 2.6 (Minimizers converge to minimizers under \( \Gamma \)-convergence). Assume \( F_n \) \( \Gamma \)-converges to \( F \) in the sense of Definition 2.7. If for every \( n \), \( x_n \) minimizes \( F_n \), and if the sequence \( \{ x_n \} \) converges to some \( x \) in \( X \), then \( x \) minimizes \( F \), and moreover, \( \lim_{n \to +\infty} \min_X F_n = \min_X F \).

Proof. Let \( y \in X \). By the \( \Gamma \)-lim sup inequality, there is a recovery sequence \( \{ y_n \} \) converging to \( y \) such that \( F(y) \geq \limsup_{n \to +\infty} F_n(y_n) \). By minimality of \( x_n \), we have \( F_n(y_n) \geq F_n(x_n) \) for all \( n \) and by the \( \Gamma \)-lim inf inequality it follows that \( \liminf_{n \to +\infty} F_n(x_n) \geq F(x) \), hence \( F(y) \geq F(x) \). Since this is true for every \( y \) in \( X \), it proves that \( x \) is a minimizer of \( F \). The relation \( \lim_{n \to +\infty} \min F_n = \min F \) follows from the previous chain of inequalities applied with \( y = x \). \( \square \)

Remark 2.7. An additional compactness assumption as in Remark 2.2 ensures that if \( \{ \min F_n \} \) is bounded then a sequence \( \{ x_n \} \) of minimizers has a limit, up to extraction. That limit must then be a minimizer of \( F \). If moreover it happens that \( F \) has a unique minimizer, then the whole sequence \( \{ x_n \} \) must converge to it.
2.2 Γ-convergence of the Coulomb gas Hamiltonian

The example of Γ-convergence of interest for us here is that of the sequence of functions \( \{ \frac{1}{n^2} H_n \} \), defined as in (2.1). The space \( \mathcal{P}(\mathbb{R}^d) \) of Borel probability measures on \( \mathbb{R}^d \) endowed with the topology of weak convergence (i.e. that of the dual of bounded continuous functions in \( \mathbb{R}^d \)), which is metrizable, will play the role of the metric space \( X \) above. We may view \( H_n \) as being defined on \( \mathcal{P}(\mathbb{R}^d) \) through the map

\[
\left\{ \begin{array}{l}
(\mathbb{R}^d)^n \\ (x_1, \ldots, x_n)
\end{array} \right\} \rightarrow \mathcal{P}(\mathbb{R}^d),
\]

which associates to any configuration of \( n \) points the probability measure \( \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i} \), also called the empirical measure, or spectral measure in the context of random matrices. More precisely, for any \( \mu \in \mathcal{P}(\mathbb{R}^d) \), we let \( H_n(\mu) \) be:

\[
H_n(\mu) = \left\{ \begin{array}{ll}
H_n(x_1, \ldots, x_n) & \text{if } \mu \text{ is of the form } \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i} \\
+\infty & \text{otherwise.}
\end{array} \right.
\]

The first main result that we prove here is that the sequence \( \{ \frac{1}{n^2} H_n \} \), has an explicit Γ-limit. It should not be surprising that we first need to divide by \( n^2 \) in order to get a limit, since we have seen that all the terms in \( H_n \) are expected to be proportional to \( n^2 \).

**Proposition 2.8** (Γ-convergence of \( \frac{1}{n^2} H_n \)). Assume \( V \) is continuous and bounded below. The sequence \( \{ \frac{1}{n^2} H_n \} \) of functions (defined on \( \mathcal{P}(\mathbb{R}^d) \) as above) Γ-converges as \( n \to +\infty \), with respect to the weak convergence of probability measures, to the function \( I : \mathcal{P}(\mathbb{R}^d) \to (-\infty, +\infty] \) defined by:

\[
I(\mu) := \iint_{\mathbb{R}^d \times \mathbb{R}^d} g(x - y) d\mu(x) d\mu(y) + \int_{\mathbb{R}^d} V(x) d\mu(x). \tag{2.7}
\]

Note that \( I(\mu) \) is simply a continuous version of the discrete Hamiltonian \( H_n \) defined over all \( \mathcal{P}(\mathbb{R}^d) \), which may also take the value \( +\infty \). From the point of view of statistical mechanics, \( I \) is the “mean-field” limit energy of \( H_n \), while we will see below in Section 2.6 that from the point of view of probability, \( I \) also plays the role of a rate function.

In the next sections, we focus on the analysis of \( I \) and its minimization. This analysis will provide ingredients for the proof of Proposition 2.8 which is postponed to Section 2.4.
2.3 Minimizing the mean-field energy via potential theory

In this section, we focus on the study of $I$ defined in (2.7), and the associated minimization problem of finding

$$
\min_{\mu \in \mathcal{P}(\mathbb{R}^d)} \int_{\mathbb{R}^d \times \mathbb{R}^d} g(x - y)d\mu(x)d\mu(y) + \int_{\mathbb{R}^d} V(x)d\mu(x). \tag{2.8}
$$

This turns out to be a classical electrostatic problem, that of finding the equilibrium distribution of charges in a capacitor with an external potential, also called the “capacitance problem,” historically studied by Gauss, and rigorously solved by Frostman [Fro]. It thus is a fundamental question in potential theory, which itself grew out of the study of the electrostatic or gravitational potential. One may see e.g. [AdHed, Doob, SaTo] and references therein. The case of $d = 2$ and $g(x) = -\log |x|$ is precisely treated in [SaTo, Chap. 1]. Higher dimensional and more general singular interaction potentials are treated in [CGZ]. The general case is not more difficult.

We start with

**Lemma 2.9.** The functional $I$ is strictly convex on $\mathcal{P}(\mathbb{R}^d)$.

**Proof.** Since $\mu \mapsto \int V d\mu$ is linear, it suffices to notice that the quadratic function $f \mapsto \int g(x - y)df(x)df(y)$, defined over all (signed) Radon measures, is positive. This is true in dimension $\geq 3$ (cf. [LiLo, Theorem 9.8]), or when $g = -\log$ in dimension 2 (for a proof cf. [RSY2, Lemma 3.2], which itself relies on [SaTo, Chap. 1, Lemma 1.8].)

Note that less restrictive assumptions, such as $g > 0$ or $\hat{g} > 0$, where $\hat{g}$ stands for the Fourier transform, would also suffice.

As a consequence, there is a unique (if any) minimizer to (2.8). In potential theory it is called the equilibrium measure or the Frostman equilibrium measure, or sometimes the extremal measure.

The question of the existence of a minimizer is a bit more delicate. In order to show it, we start by making the following assumptions on the potential $V$:

(A1) $V$ is l.s.c. and bounded below.

(A2) (growth assumption)

$$
\lim_{|x| \to +\infty} \left( \frac{V(x)}{2} + g(x) \right) = +\infty
$$

The first condition is there to ensure the lower semi-continuity of $I$ and that $\inf I > -\infty$, the second is made to ensure that $I$ is coercive. Of course, in the Coulomb case with $d \geq 3$, condition (A2) is equivalent to the condition that $V$ tends to $+\infty$ at infinity.
Lemma 2.10. Assume (A1) and (A2) are satisfied, and let \( \{\mu_n\}_n \) be a sequence in \( \mathcal{P}(\mathbb{R}^d) \) such that \( \{I(\mu_n)\}_n \) is bounded. Then, up to extraction of a subsequence \( \mu_n \) converges to some \( \mu \in \mathcal{P}(\mathbb{R}^d) \) in the weak sense of probabilities, and

\[
\liminf_{n \to \infty} I(\mu_n) \geq I(\mu).
\]

Proof. Assume that \( I(\mu_n) \leq C_1 \) for each \( n \). Given any constant \( C_2 > 0 \) there exists a compact set \( K \subset \mathbb{R}^d \) such that

\[
\min_{(K \times K)} \left[ g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \right] > C_2. \tag{2.9}
\]

Indeed, it suffices to show that \( g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \to \infty \) as \( x \to \infty \) or \( y \to \infty \). To check this, one may separate the cases \( d = 1, 2 \) and \( d \geq 3 \). In the latter case, the Coulomb kernel \( g \) is positive and \( V \) is bounded below, so that \( g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \), which is greater than \( \frac{V}{2}(x) + \frac{V}{2}(y) \), can be made arbitrarily large if \( x \) or \( y \) gets large by assumption (A2). When \( d = 1, 2 \), since

\[
g(x - y) = -\log|y| \geq -\log 2 - \log \max(|x|, |y|),
\]

we also have from assumptions (A1) and (A2) that \( \frac{1}{2}(V(x) + V(y)) + g(x - y) \) is arbitrarily large if \( |x| \) or \( |y| \) is large enough.

In addition, by (A1), (A2), we have in all cases that \( g(x - y) + \frac{1}{2}V(x) + \frac{1}{2}V(y) \) is bounded below on \( \mathbb{R}^d \times \mathbb{R}^d \) by a constant, say \( -C_3 \), with \( C_3 > 0 \). Rewriting then \( I \) as

\[
I(\mu) = \iint_{\mathbb{R}^d} \left[ g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \right] d\mu(x)d\mu(y), \tag{2.10}
\]

the relation (2.9) and our assumption on \( \mu_n \) imply that

\[
C_1 \geq I(\mu_n) \geq -C_3 + C_2(\mu_n \otimes \mu_n)((K \times K)^c) \geq -C_3 + C_2\mu_n(K^c).
\]

Since \( C_2 \) can be made arbitrarily large, \( \mu_n(K^c) \) can be made arbitrarily small, which means precisely that \( \{\mu_n\}_n \) is a tight sequence of probability measures. By Prokhorov’s theorem, it thus has a convergent subsequence (still denoted \( \{\mu_n\}_n \)), which converges to some probability \( \mu \). For any \( n \) and any \( M > 0 \), we may then write

\[
\iint g(x - y)d\mu_n(x)d\mu_n(y) \geq \iint (g(x - y) \wedge M)d\mu_n(x)d\mu_n(y) \tag{2.11}
\]

where \( \wedge \) denotes the minimum of two numbers. For each given \( M \), the integrand in the right-hand side is continuous, and thus the weak convergence of \( \mu_n \) to \( \mu \), which implies the weak convergence of \( \mu_n \otimes \mu_n \) to \( \mu \otimes \mu \), yields

\[
\liminf_{n \to +\infty} \iint g(x - y)d\mu_n(x)d\mu_n(y) \geq \iint (g(x - y) \wedge M)d\mu(x)d\mu(y).
\]

A monotone convergence theorem argument allows one to let \( M \to +\infty \), and using assumption (A1) for the weak l.s.c of the potential part of the functional, we conclude that

\[
\liminf_{n \to +\infty} I(\mu_n) \geq I(\mu). \tag{2.12}
\]
We have seen above that \( \inf I > -\infty \) (indeed the integrand in \( I \) is bounded below thanks to assumption (A1)). The next question is to see whether \( \inf I < +\infty \), i.e. that there exist probabilities with finite \( I \)'s. This is directly related to the notion of (electrostatic, Bessel, or logarithmic) capacity, whose definition we now give. One may find it \( [\text{SaTo}, \text{EvGar}, \text{AdHed}] \) or \( [\text{LiLo}, \text{Sec. 11.15}] \), the formulations differ a bit but are essentially equivalent. It is usually not formulated this way in dimension 1 but it can be extended to that case with our choice of \( g = -\log \) without trouble.

**Definition 2.11** (Capacity of a set). We define the capacity of a compact set \( K \subset \mathbb{R}^d \) by

\[
\text{cap}(K) := \Phi \left( \inf_{\mu \in \mathcal{P}(K)} \int_{\mathbb{R}^d} g(x - y)d\mu(x)d\mu(y) \right),
\]

with \( \Phi(t) = e^{-t} \) if \( d = 1, 2 \) and \( \Phi(t) = t^{-\frac{d}{2}} \) if \( d \geq 3 \), and where \( \mathcal{P}(K) \) denotes the set of probability measures supported in \( K \). Here the inf can be \( +\infty \) if there exists no probability measure \( \mu \in \mathcal{P}(K) \) such that \( \int_{\mathbb{R}^d} g(x - y)d\mu(x)d\mu(y) < +\infty \). For a general set \( E \), we define \( \text{cap}(E) \) as the supremum of \( \text{cap}(K) \) over the compact sets \( K \) included in \( E \). It is easy to check that capacity is increasing with respect to the inclusion of sets.

A basic fact is that a set of zero capacity also has zero Lebesgue measure (see the references above).

**Lemma 2.12.** If \( \text{cap}(E) = 0 \), then \( |E| = 0 \).

In fact \( \text{cap}(E) = 0 \) is stronger than \( |E| = 0 \), it implies for example that the perimeter of \( E \) is also 0. A property is said to hold “quasi-everywhere” (q.e.) if it holds everywhere except on a set of capacity zero. By the preceding lemma a property that holds q.e. also holds Lebesgue-almost everywhere (a.e.), whereas the converse is, in general, not true.

For the sake of generality, it is interesting to consider potential \( V \)'s which can take the value \( +\infty \) (this is the same as imposing the constraint that the probability measures only charge a specific set, the set where \( V \) is finite). We then need to place a third assumption

(A3) \( \{ x \in \mathbb{R}^d | V(x) < +\infty \} \) has positive capacity.

**Lemma 2.13.** Under assumptions (A1) – (A3), we have \( \inf I < +\infty \).

**Proof.** Let us define for any \( \varepsilon > 0 \) the set \( \Sigma_{\varepsilon} = \{ x \ | \ V(x) \leq \frac{1}{\varepsilon} \} \). Since \( V \) is l.s.c the sets \( \Sigma_{\varepsilon} \) are closed, and it is easy to see that assumption (A2) implies that they are also bounded, since \( V(x) \) goes to \( +\infty \) when \( |x| \to +\infty \).

The capacity of \( \Sigma_0 = \{ x \in \mathbb{R}^d | V(x) < +\infty \} \) is positive by assumption. It is easily seen that the sets \( \{ \Sigma_{\varepsilon} \}_{\varepsilon > 0} \) form a decreasing family of compact sets with \( \bigcup_{\varepsilon > 0} \Sigma_{\varepsilon} = \Sigma_0 \), and by definition (see Definition 2.11 or the references given above)
2.3 Minimizing the mean-field energy via potential theory

the capacity of $\Sigma_0$ is the supremum of capacities of compact sets included in $\Sigma_0$. Hence we have that cap($\Sigma_\varepsilon$) is positive for $\varepsilon$ small enough. Then by definition there exists a probability measure $\mu_\varepsilon$ supported in $\Sigma_\varepsilon$ such that

$$\int g(x-y)d\mu_\varepsilon(x)d\mu_\varepsilon(y) < +\infty. \quad (2.14)$$

Of course, we also have $\int Vd\mu_\varepsilon < +\infty$ by definition of $\Sigma_\varepsilon$. Hence $I(\mu_\varepsilon) < +\infty$, in particular $\inf I < +\infty$.

We may now give the main existence result, together with the characterization of the minimizer.

**Theorem 2.1** (Frostman [Fro], existence and characterization of the equilibrium measure). Under the assumptions (A1)-(A2)-(A3), the minimum of $I$ over $\mathcal{P}(\mathbb{R}^d)$ exists, is finite and is achieved by a unique $\mu_0$, which has a compact support of positive capacity. In addition $\mu_0$ is uniquely characterized by the fact that

$$\begin{cases}
    h^{\mu_0} + \frac{V}{2} \geq c \quad &\text{q.e. in } \mathbb{R}^d \\
    h^{\mu_0} + \frac{V}{2} = c \quad &\text{q.e. in the support of } \mu_0
\end{cases} \quad (2.15)$$

where

$$h^{\mu_0}(x) := \int_{\mathbb{R}^d} g(x-y)d\mu_0(y) \quad (2.16)$$

is the electrostatic potential generated by $\mu_0$; and then the constant $c$ must be

$$c = I(\mu_0) - \frac{1}{2} \int_{\mathbb{R}^d} V(x)d\mu_0(x). \quad (2.17)$$

**Remark 2.14.** Note that by (2.3), in dimension $d \geq 2$, the function $h^{\mu_0}$ solves:

$$-\Delta h^{\mu_0} = c_d \mu_0.$$

**Example 2.15** (Capacity of a compact set). Let $K$ be a compact set of positive capacity, and let $V = 0$ in $K$ and $V = +\infty$ in $K^c$. In that case the minimization of $I$ is the same as the computation of the capacity of $K$ as in (2.13). The support of the equilibrium measure $\mu_0$ is contained in $K$, and the associated Euler-Lagrange equation (2.15) states that the electrostatic potential (if in dimension $d \geq 2$) $h^{\mu_0}$ is constant q.e. on the support of $\mu_0$ (a well-known result in physics). If $K$ is sufficiently regular, one can apply the Laplacian on both sides of this equality, and in view of Remark 2.14 find that $\mu_0 = 0$ q.e. in $K$, which indicates that $\mu_0$ is supported on $\partial K$. 
The leading order behavior of the Coulomb gas

Example 2.16 (C$^{1,1}$ potentials and RMT examples). In general, the relations (2.15) say that the total potential $h^{\mu_0} + \frac{1}{2}V$ is constant on the support of the charges. Moreover, in dimension $d \geq 2$, applying the Laplacian on both sides of (2.15) and using again Remark 2.14 gives that, on the interior of the support of the equilibrium measure, if $V \in C^{1,1}$,

$$c_d \mu_0 = \frac{\Delta V}{2} \tag{2.18}$$

(where $c_d$ is the constant defined in (2.4)), i.e. the density of the measure on the interior of its support is given by $\frac{\Delta V}{\Delta c_d}$. This will be proven in Proposition 2.22. For example if $V$ is quadratic, then the measure $\mu_0 = \frac{\Delta V}{\Delta c_d}$ is constant on the interior of its support. This corresponds to the important examples of the Hamiltonians that arise in random matrix theory, more precisely:

- in dimension $d = 2$, for $V(x) = |x|^2$, one may check that $\mu_0 = \frac{1}{\pi} 1_{B_1}$, where $1$ denotes a characteristic function and $B_1$ is the unit ball, i.e. the equilibrium measure is the normalized Lebesgue measure on the unit disk (by uniqueness, $\mu_0$ should be radially symmetric, and the combination of (2.18) with the constraint of being a probability measure imposes the support to be $B_1$). This is known as the circle law for the Ginibre ensemble in the context of Random Matrix Theory (RMT). Its derivation (which we will see in Section 2.6 below) is attributed to Ginibre, Mehta, an unpublished paper of Silverstein and Girko [Gi].

- in dimension $d \geq 3$, the same holds, i.e. for $V(x) = |x|^2$ we have $\mu_0 = \frac{d}{c_d} 1_{B_{(d-2)/d}}$ by the same reasoning.

- in dimension $d = 1$, with $g = -\log |\cdot|$ and $V(x) = x^2$, the equilibrium measure is $\mu_0(x) = \frac{1}{\pi} \sqrt{4-x^2} 1_{|x|\leq 2}$, which corresponds in the context of RMT (GUE and GOE ensembles) to Wigner’s semi-circle law, cf. [Wi,Me].

We now turn to the proof of Theorem 2.1, adapted from [SaTo, Chap. 1].

Proof. The existence of a minimizer $\mu_0$ follows directly from Lemmas 2.10 and 2.13, its uniqueness from Lemma 2.9. It remains to show that $\mu_0$ has compact support of finite capacity and that (2.15) holds.

Step 1. We prove that $\mu_0$ has compact support. Using (2.9), we may find a compact set $K$ such that $g(x - y) + \frac{1}{2}V(x) + \frac{1}{2}V(y) \geq I(\mu_0) + 1$ outside of $K \times K$.

Assume that $\mu_0$ has mass outside $K$, i.e. assume $\mu_0(K) < 1$, and define the new probability measure

$$\tilde{\mu} := \frac{(\mu_0|_K)}{\mu_0(K)} \tag{2.19}$$

We want to show that $\tilde{\mu}$ has less or equal energy $I(\tilde{\mu})$ than $\mu_0$, in order to get a
contradiction. One may compute $I(\mu_0)$ in the following way:

$$I(\mu_0) = \int_{K \times K} \left[ g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \right] d\mu_0(x) d\mu_0(y)$$

$$+ \int_{(K \times K)^c} \left[ g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \right] d\mu_0(x) d\mu_0(y)$$

$$\geq \mu_0(K)^2 I(\tilde{\mu}) + (1 - \mu_0(K))^2 \min_{(K \times K)^c} \left[ g(x - y) + \frac{V}{2}(x) + \frac{V}{2}(y) \right]. \quad (2.20)$$

By choice of $K$, and since we assumed $\mu_0(K) < 1$, this implies that

$$I(\mu_0) \geq \mu_0(K)^2 I(\tilde{\mu}) + (1 - \mu_0(K))^2 (I(\mu_0) + 1) \quad (2.21)$$

and thus

$$I(\tilde{\mu}) \leq I(\mu_0) + \frac{\mu_0(K)^2 - 1}{\mu_0(K)^2} < I(\mu_0),$$

a contradiction with the minimality of $\mu_0$. We thus conclude that $\mu_0$ has compact support. The fact that the support of $\mu_0$ has positive capacity is an immediate consequence of the fact that $I(\mu_0) < \infty$ and the definition of capacity.

**Step 2.** We turn to the proof of the Euler-Lagrange equations (2.15). For this, we use the “method of variations” which consists in continuously deforming $\mu_0$ into other admissible probability measures.

Let $\nu$ in $\mathcal{P}(\mathbb{R}^d)$ such that $I(\nu) < +\infty$, and consider the probability measure $(1 - t)\mu_0 + t\nu$ for all $t \in [0, 1]$. Since $\mu_0$ minimizes $I$, we have

$$I((1 - t)\mu_0 + t\nu) \geq I(\mu_0), \quad \text{for all } t \in [0, 1]. \quad (2.22)$$

By letting $t \to 0^+$ and keeping only the first order terms in $t$, one obtains the “functional derivative” of $I$ at $\mu_0$. More precisely, writing

$$\int g(x - y)d((1 - t)\mu_0 + t\nu)(x)d((1 - t)\mu_0 + t\nu)(y) + \int V(x)d((1 - t)\mu_0 + t\nu)(x) \geq I(\mu_0) \quad (2.23)$$

one easily gets that

$$I(\mu_0) + t \left[ \int g(x - y)d(\nu - \mu_0)(x)d\mu_0(y) + \int g(x - y)d(\nu - \mu_0)(y)d\mu_0(x) \right. \quad (2.24)$$

$$\left. + \int V(x)d(\nu - \mu_0)(x) \right] + O(t^2) \geq I(\mu_0).$$

Here, we may cancel the identical order 0 terms $I(\mu_0)$ on both sides, and note that in view of (2.16) the expression between brackets can be rewritten as $2 \int h^{\mu_0}(x)d(\nu - \mu_0)(x) + \int V(x)d(\nu - \mu_0)(x)$. Next, dividing the inequality by $t > 0$, and letting $t \to 0^+$, it appears that for all $\nu$ in $\mathcal{P}(\mathbb{R}^d)$ such that $I(\nu) < +\infty$, the following inequality holds:

$$2 \int h^{\mu_0}(x)d(\nu - \mu_0)(x) + \int V(x)d(\nu - \mu_0)(x) \geq 0 \quad (2.25)$$
or equivalently
\[ \int \left( h^{\mu_0} + \frac{V}{2} \right)(x) d\nu(x) \geq \int \left( h^{\mu_0} + \frac{V}{2} \right)(x) d\mu_0(x). \]  
\( (2.26) \)

Defining the constant \( c \) by
\[ c = I(\mu_0) - \frac{1}{2} \int V d\mu_0 = \iint g(x - y) d\mu_0(x) d\mu_0(y) + \frac{1}{2} \int V d\mu_0 \]
\[ = \int \left( h^{\mu_0} + \frac{V}{2} \right) d\mu_0, \]
\( (2.27) \)

\( (2.26) \) asserts that
\[ \int \left( h^{\mu_0} + \frac{V}{2} \right) d\nu \geq c \]
\( (2.28) \)

for all probability measures \( \nu \) on \( \mathbb{R}^d \) such that \( I(\nu) < +\infty \). Note that at this point, if we relax the condition \( I(\nu) < +\infty \), then choosing \( \nu \) to be a Dirac mass when applying \( (2.28) \) would yield
\[ h^{\mu_0} + \frac{V}{2} \geq c \]
\( (2.29) \)

pointwise. However, Dirac masses have infinite energy \( I \), and we can only prove that \( (2.29) \) holds quasi-everywhere, which we do now.

Assume not, then there exists a set \( K \) of positive capacity such that \( (2.29) \) is false on \( K \), and by definition of the capacity of \( K \) as as supremum of capacities over compact sets included in \( K \), we may in fact suppose that \( K \) is compact. By definition, this means that there is a probability measure \( \nu \) supported in \( K \) such that
\[ \iint g(x - y) d\nu(x) d\nu(y) < +\infty. \]
\( (2.30) \)

Let us observe that \( -h^{\mu_0} \) is bounded above on any compact set (this is clear in dimension \( d \geq 3 \) because the Coulomb kernel is positive and so is \( h^{\mu_0} \), and can be easily checked in \( d = 1, 2 \) because log is bounded above on any compact set and \( \mu_0 \) has compact support). By assumption, equation \( (2.29) \) is false on \( K \), that is \( V < 2c - 2h^{\mu_0} \) on \( K \). Integrating this inequality against \( \nu \) gives
\[ \int V d\nu = \int_K V d\nu < \int_K (2c - 2h^{\mu_0}) d\nu < +\infty \]
\( (2.31) \)

which, combined with \( (2.30) \) ensures that \( I(\nu) \) is finite. But then \( (2.31) \) contradicts \( (2.28) \). We thus have shown that
\[ h^{\mu_0} + \frac{V}{2} \geq c \text{ q.e.} \]
\( (2.32) \)

which is the first of the relations \( (2.15) \).
2.3 Minimizing the mean-field energy via potential theory

For the second one, let us denote by $E$ the set where the previous inequality (2.32) fails. We know that $E$ has zero capacity, but since $\mu_0$ satisfies $I(\mu_0) < +\infty$, it does not charge sets of zero capacity (otherwise one could restrict $\mu_0$ on such a set, normalize its mass to 1 and get a contradiction with the definition of a zero capacity set). Hence we have

$$h^{\mu_0} + \frac{V}{2} \geq c \quad \mu_0 \text{-a.e.} \quad (2.33)$$

Integrating this relation against $\mu_0$ yields

$$\hat{\int} (h^{\mu_0} + \frac{V}{2}) d\mu_0 \geq c, \quad (2.34)$$

but in view of (2.27) this implies that equality must hold in (2.33) $\mu_0$-almost everywhere. This establishes the second Euler-Lagrange equation.

**Step 3.** We show that the relations (2.15) uniquely characterize the minimizer of $I$. Assume that $\mu$ is another probability solving (2.15) with some constant $c'$, and set, for $t \in (0, 1)$, $\mu_t := t\mu + (1 - t)\mu_0$, hence $h^{\mu_t} = t h^{\mu} + (1 - t) h^{\mu_0}$. We have

$$I(\mu_t) = \int (t h^{\mu}(x) + (1 - t) h^{\mu_0}(x) + V(x)) d\mu_t(x)$$

$$= \frac{t}{2} \int (2 h^{\mu}(x) + V(x)) d\mu_t(x)$$

$$+ \frac{(1 - t)}{2} \int (2 h^{\mu_0}(x) + V(x)) d\mu_t(x) + \frac{1}{2} \int V(x) d\mu_t(x).$$

By assumption, $h^{\mu} + \frac{V}{2} \geq c'$ and $h^{\mu_0} + \frac{V}{2} \geq c$ almost everywhere. We hence get that

$$I(\mu_t) \geq tc' + (1 - t)c + \frac{1}{2} \int V(x) \left( t d\mu(x) + (1 - t) d\mu_0(x) \right)$$

$$= t \left( c' + \frac{1}{2} \int V d\mu \right) + (1 - t) \left( c + \frac{1}{2} \int V d\mu_0 \right). \quad (2.35)$$

On the other hand, integrating the second Euler-Lagrange equation in (2.15) for both $\mu$ and $\mu_0$, with respect to $\mu$ and $\mu_0$ respectively, yields, after rearranging terms,

$$I(\mu) = c' + \frac{1}{2} \int V d\mu \quad \text{and} \quad I(\mu_0) = c + \frac{1}{2} \int V d\mu_0.$$

Hence $I(\mu_t) \geq t I(\mu) + (1 - t) I(\mu_0)$, which is impossible by strict convexity of $I$ unless $\mu = \mu_0$. This proves that the two measures $\mu$ and $\mu_0$ must coincide. \(\Box\)

**Remark 2.17.** In all this section, we did not use much all the particulars of the Coulomb kernel. The theorem still holds for a much more general class of $g$’s, say $g$ positive, monotone radial and satisfying $\iint g(x - y) \, dx \, dy < \infty$. 

For the second one, let us denote by $E$ the set where the previous inequality (2.32) fails. We know that $E$ has zero capacity, but since $\mu_0$ satisfies $I(\mu_0) < +\infty$, it does not charge sets of zero capacity (otherwise one could restrict $\mu_0$ on such a set, normalize its mass to 1 and get a contradiction with the definition of a zero capacity set). Hence we have

$$h^{\mu_0} + \frac{V}{2} \geq c \quad \mu_0 \text{-a.e.} \quad (2.33)$$

Integrating this relation against $\mu_0$ yields

$$\hat{\int} (h^{\mu_0} + \frac{V}{2}) d\mu_0 \geq c, \quad (2.34)$$

but in view of (2.27) this implies that equality must hold in (2.33) $\mu_0$-almost everywhere. This establishes the second Euler-Lagrange equation.

**Step 3.** We show that the relations (2.15) uniquely characterize the minimizer of $I$. Assume that $\mu$ is another probability solving (2.15) with some constant $c'$, and set, for $t \in (0, 1)$, $\mu_t := t\mu + (1 - t)\mu_0$, hence $h^{\mu_t} = t h^{\mu} + (1 - t) h^{\mu_0}$. We have

$$I(\mu_t) = \int (t h^{\mu}(x) + (1 - t) h^{\mu_0}(x) + V(x)) d\mu_t(x)$$

$$= \frac{t}{2} \int (2 h^{\mu}(x) + V(x)) d\mu_t(x)$$

$$+ \frac{(1 - t)}{2} \int (2 h^{\mu_0}(x) + V(x)) d\mu_t(x) + \frac{1}{2} \int V(x) d\mu_t(x).$$

By assumption, $h^{\mu} + \frac{V}{2} \geq c'$ and $h^{\mu_0} + \frac{V}{2} \geq c$ almost everywhere. We hence get that

$$I(\mu_t) \geq tc' + (1 - t)c + \frac{1}{2} \int V(x) \left( t d\mu(x) + (1 - t) d\mu_0(x) \right)$$

$$= t \left( c' + \frac{1}{2} \int V d\mu \right) + (1 - t) \left( c + \frac{1}{2} \int V d\mu_0 \right). \quad (2.35)$$

On the other hand, integrating the second Euler-Lagrange equation in (2.15) for both $\mu$ and $\mu_0$, with respect to $\mu$ and $\mu_0$ respectively, yields, after rearranging terms,

$$I(\mu) = c' + \frac{1}{2} \int V d\mu \quad \text{and} \quad I(\mu_0) = c + \frac{1}{2} \int V d\mu_0.$$

Hence $I(\mu_t) \geq t I(\mu) + (1 - t) I(\mu_0)$, which is impossible by strict convexity of $I$ unless $\mu = \mu_0$. This proves that the two measures $\mu$ and $\mu_0$ must coincide. \(\Box\)

**Remark 2.17.** In all this section, we did not use much all the particulars of the Coulomb kernel. The theorem still holds for a much more general class of $g$’s, say $g$ positive, monotone radial and satisfying $\iint g(x - y) \, dx \, dy < \infty$. 

Definition 2.18. From now on, we denote by \( \zeta \) the function

\[
\zeta = h^{\mu_0} + \frac{V}{2} - c. \tag{2.36}
\]

We note that in view of (2.15), \( \zeta \geq 0 \) a.e. and \( \zeta = 0 \) \( \mu_0 \)-a.e.

2.4 Proof of the \( \Gamma \)-convergence and consequences for minimizers of the Hamiltonian: the mean-field limit

We now proof Proposition 2.8. The proof uses the same ingredients as the proof of the existence of a minimizer of \( I \) in the previous section. A statement and a proof with \( \Gamma \)-convergence in dimension 2 for \( V \) quadratic appeared in [SS4, Proposition 11.1]. It is not difficult to adapt them to higher dimensions and more general potentials. Similar arguments are also found in the large deviations proofs of [BG, BZ, CGZ].

In what follows, when considering sequences of configurations \((x_1, \ldots, x_n)\) we will make the slight abuse of notation that consists in neglecting the dependency of the points \((x_1, \ldots, x_n)\) on \( n \), while one should formally write \((x_{1,n}, \ldots, x_{n,n})\).

Proof of Proposition 2.8. In the following, we denote the diagonal of \( \mathbb{R}^d \times \mathbb{R}^d \) by \( \Delta \) and its complement by \( \Delta^c \).

Step 1. We first need to prove that if \( \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i} \rightharpoonup \mu \in \mathcal{P}(\mathbb{R}^d) \), then

\[
\liminf_{n \to +\infty} \frac{1}{n^2} H_n(x_1, \ldots, x_n) \geq I(\mu). \tag{2.37}
\]

Letting \( \mu_n \) denote the empirical measure \( \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i} \), we may write

\[
\frac{1}{n^2} H_n(\mu_n) = \iint_{\Delta^c} g(x-y) d\mu_n(x) d\mu_n(y) + \int V(x) d\mu_n(x). \tag{2.38}
\]

The last term in the right-hand side is harmless: since \( V \) is assumed to be continuous and bounded below, it is lower-semicontinuous and bounded below and since the sequence \( \{\mu_n\} \) converges weakly to \( \mu \), we have:

\[
\liminf_{n \to +\infty} \int V d\mu_n \geq \int V d\mu. \tag{2.39}
\]

In order to treat the first term in the right-hand side of (2.37), following [SaTo, Chap. 1], let us truncate the singularity of \( g \) by writing:

\[
\iint_{\Delta^c} g(x-y) d\mu_n(x) d\mu_n(y) \geq \iint (g(x-y) \wedge M) d\mu_n(x) d\mu_n(y) - \frac{M}{n}. \tag{2.39}
\]
where $M > 0$ and $\wedge$ still denotes the minimum of two numbers. Indeed one has $\mu_n \otimes \mu_n(\triangle) = \frac{1}{n}$, as soon as the points of the configuration $(x_1, \ldots, x_n)$ are simple (i.e. $x_i \neq x_j$ for $i \neq j$). The function $(x, y) \mapsto g(x - y) \wedge M$ is continuous, and by taking the limit of (2.39) as $n \to +\infty$ one gets, by weak convergence of $\mu_n$ to $\mu$ (hence of $\mu_n \otimes \mu_n$ to $\mu \otimes \mu$) that for every $M > 0$:

$$\liminf_{n \to +\infty} \int_{\triangle} g(x - y)d\mu_n(x)d\mu_n(y) \geq \int (g(x - y) \wedge M)d\mu(x)d\mu(y). \quad (2.40)$$

By the monotone convergence theorem, the (possibly infinite) limit of the right-hand side as $M \to +\infty$ exists and equals $\int g(x - y)d\mu(x)d\mu(y)$. Combining with (2.38) and (2.39), this concludes the proof of the $\Gamma$-lim inf convergence. Let us note that for this part, we really only needed to know that $V$ is lower semi-continuous and bounded below.

**Step 2.** We now need to construct a recovery sequence for each measure $\mu$ in $\mathcal{P}(\mathbb{R}^d)$. First, we show that we can reduce to measures which are in $L^\infty(\mathbb{R}^d)$, supported in a cube $K$ and such that the density $\mu(x)$ is bounded below by $\alpha > 0$ in $K$. Let $\mu$ be an arbitrary measure in $\mathcal{P}(\mathbb{R}^d)$ such that $I(\mu) < +\infty$. Given $\alpha > 0$, by tightness of $\mu$, we may truncate it outside of a compact set which contains all its mass but $\alpha$. Making the compact set larger if necessary, and normalizing the truncated measure to make it a probability, the argument of Step 1 of the proof of Theorem 2.1 shows that this decreases $I$. In other words we have a family $\mu_\alpha$ with $\mu_\alpha \to \mu$ and $\limsup_{\alpha \to 0} I(\mu_\alpha) \leq I(\mu)$. Thus, by a diagonal argument, it suffices to prove our statement for probability measures $\mu$ which have compact support.

Let us next consider such a probability $\mu$. Convolving $\mu$ with smooth mollifiers $\chi_\eta$ (positive of integral 1), we may approximate $\mu$ by smooth $\mu_\eta$, these converge to $\mu$ in the weak sense of probabilities, as $\eta \to 0$. Let us denote

$$\Phi(\mu) = \int g(x - y) d\mu(x) d\mu(y).$$

As seen in the proof of Lemma 2.9, the function $\Phi$ is strictly convex. Writing $\tau_y \mu$ for the translate of $\mu$ by $y$, we deduce, using Jensen’s inequality, that

$$\Phi(\chi_\eta \star \mu) = \Phi \left( \int \tau_y \mu \chi_\eta(y) dy \right) \leq \int \chi_\eta(y) \Phi(\tau_y \mu) dy.$$

Since $\Phi$ is translation-invariant, we have $\Phi(\tau_y \mu) = \Phi(\mu)$ and thus we have obtained $\Phi(\mu_\eta) \leq \Phi(\mu)$. On the other hand, $\lim_{\eta \to 0} \int Vd(\mu_\eta) = \int Vd\mu$ since $V$ is assumed to be continuous and $\mu_\eta \to \mu$ and they all are supported in the same compact set. We have thus established that $\limsup_{\eta \to 0} I(\mu_\eta) \leq I(\mu)$. Thus, by a diagonal argument, it suffices to prove our statement for probability measures $\mu$ which have a smooth density and compact support.

Let us next consider such a probability measure $\mu$. We may find a cube $K$ that contains its support and consider the probability measure $\mu_\alpha = \frac{\mu + \alpha \chi_K}{1 + \alpha |K|}$. It is supported in the cube $K$, has an $L^\infty$ density which is bounded below by $\alpha$ in $K$, 

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as desired, and \( \mu_n \to \mu \) in the weak sense of probabilities, but also in \( L^\infty \). It is easy to deduce from this fact that since \( \iint g(x-y) \, dx \, dy < +\infty \) (cf. (2.2)) and \( V \) is continuous, we have \( I(\mu_n) \to I(\mu) \). Again, a diagonal argument allows us to reduce to proving the desired \( \Gamma \)-limsup statement for such measures.

**Step 3.** Henceforth we assume that \( \mu \) is in \( L^\infty(\mathbb{R}^d) \), supported in a cube \( K \) and such that the density \( \mu(x) \) is bounded below by \( \alpha > 0 \) in \( K \). Since we are going to construct a configuration of \( n \) points in a cube \( K \), the typical lengthscale of the distance between two points is \( \frac{1}{n^{1/d}} \). Let us then choose a sequence \( c_n \) such that \( \frac{1}{n^{1/d}} \ll c_n \ll 1 \) as \( n \to +\infty \), and for each \( n \), split \( K \) into cubes \( K_k \) (depending on \( n \)) whose sidelength is in \([c_n, 2c_n]\), cf. Fig. 2.1.

**Claim 1.** We may place \( n_k = [n\mu(K_k)] \pm 0 \), 1 points in \( K_k \) (where \( [\cdot] \) denotes the integer part), with \( \sum_k n_k = n \), and such that the resulting sets of points \( \{x_i\}_{i=1}^n \) satisfy that the balls \( B(x_i, \frac{4}{n^{1/d}}) \) are disjoint, for some \( \lambda > 0 \) independent of \( n \).

It is possible to do so because the density of \( \mu \) is bounded above and below on \( K \), for a proof see e.g. [SS4, Lemma 7.4 and below]. Then a Riemann sum argument, combined with the facts that \( |n_k - n\mu(K_k)| \leq 1 \) and \( c_n \ll 1 \), easily allows to show that the measure \( \mu_n := \frac{1}{n} \sum_{i=1}^n \delta_{x_i} \) converges weakly to \( \mu \). We are then left with estimating \( \frac{1}{n^2} H_n(x_1, \ldots, x_n) \) from above.

For each \( 0 < \eta < 1 \) let us select \( \chi_\eta \) a smooth function on \( \mathbb{R}^d \), radial and such that \( \chi_\eta(x) = 0 \) if \( |x| < \frac{1}{2} \eta \) and \( \chi_\eta = 1 \) when \( |x| \geq \eta \). We may write

\[
\frac{1}{n^2} \iint \Delta_r g(x-y) \, dx \, dy \, d\mu_n(x) \, d\mu_n(y) = \frac{1}{n^2} \left( \iint \Delta_r [(1 - \chi_\eta)g] \, dx \, dy \, d\mu_n(x) \, d\mu_n(y) + \iint \chi_\eta g \, dx \, dy \, d\mu_n(x) \, d\mu_n(y) \right). \tag{2.41}
\]

The notation \( a_n \ll b_n \) means that \( a_n = o(b_n) \).
Since the function \( \chi_n g \) is bounded and continuous on the cube \( K \) where \( \mu_n \) and \( \mu \) are supported, the last term in the right-hand side converges to \( \iint [\chi_n g](x - y) d\mu(x) d\mu(y) \) by weak convergence of \( \mu_n \) to \( \mu \). We next turn to the first term in the right-hand side of (2.41), and show that there is, in fact, no problem near the diagonal because we have sufficient control on the accumulation of points.

Since \( 1 - \chi_n \) is bounded by 1 and vanishes outside \( B(0, \eta) \), we may write, by definition of \( \mu_n \) and positivity of \( g \) in \( B(0, 1) \) (which is true in all dimensions):

\[
\iint_{\Delta^c} [(1 - \chi_n)g](x - y) d\mu_n(x) d\mu_n(y) \leq \frac{1}{n^2} \sum_{i \neq j, \ |x_i - x_j| < \eta} g(x_i - x_j). \tag{2.42}
\]

Claim 2. For all \( x, y \in B(x_1, \frac{1}{n^{1/d}}) \times B(x_j, \frac{1}{n^{1/d}}), \ i \neq j, \) we have \( g(x_i - x_j) \leq g \left( \frac{1}{2} (x - y) \right) \).

This is due to the fact that the balls \( B(x_i, \frac{1}{n^{1/d}}) \) are disjoint from each other hence for \( i \neq j \), \( |x_i - x_j| \geq \frac{8 \lambda}{n^{1/d}} \), which implies, by the triangle inequality, that if \( x \in B(x_i, \frac{1}{n^{1/d}}) \) and \( y \in B(x_j, \frac{1}{n^{1/d}}), \) then \( |x_i - x_j| \geq \frac{1}{2} |x - y| \). But in all the cases we consider (cf. (2.2)), \( (x, y) \mapsto g(x - y) \) is a decreasing function of the distance between the two points, hence we may write:

\[
g(x_i - x_j) \leq \int_{B(x_i, \frac{1}{n^{1/d}})} \int_{B(x_j, \frac{1}{n^{1/d}})} g \left( \frac{1}{2} (x - y) \right) dxdy \\
\leq \frac{C_d}{\lambda^{2d}} \int_{B(x_i, \frac{1}{n^{1/d}})} \int_{B(x_j, \frac{1}{n^{1/d}})} g \left( \frac{1}{2} (x - y) \right) dxdy \tag{2.43}
\]

where \( C_d \) is a constant depending only on the dimension \( d \). Because the balls do not overlap, one may sum the inequalities (2.43) for \( i \neq j \) to find, with (2.42),

\[
\iint_{\Delta^c} (1 - \chi_n)g(x - y) d\mu_n(x) d\mu_n(y) \leq \frac{1}{n^2} \sum_{i \neq j, \ |x_i - x_j| < \eta} g(x_i - x_j) \\
\leq \frac{C_d}{\lambda^{2d}} \sum_{i \neq j, \ |x_i - x_j| < \eta} \int_{B(x_i, \frac{1}{n^{1/d}})} \int_{B(x_j, \frac{1}{n^{1/d}})} g \left( \frac{1}{2} (x - y) \right) dxdy \\
\leq \frac{C_d}{\lambda^{2d}} \int_{|x - y| < 2\eta} g \left( \frac{1}{2} (x - y) \right) dxdy, \tag{2.44}
\]

for \( n \) large enough. The last term in (2.44) is \( o(1) \) when \( \eta \to 0 \) because we have \( \iint g(x - y) dxdy < +\infty \). Combining with (2.41) we deduce

\[
\limsup_{n \to +\infty} \frac{1}{n^2} \iint_{\Delta^c} g(x - y) d\mu_n(x) d\mu_n(y) \leq \iint_{\mathbb{R}^d \times \mathbb{R}^d} g(x - y) d\mu(x) d\mu(y) + o_n(1). \tag{2.45}
\]

\(^2\)We will always use the notation \( \frac{1}{|U|} \int_U f \) for the average of \( f \) on \( U \), that is \( \frac{1}{|U|} \int_U f \).
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Also \( \int V d\mu_n \to \int V d\mu \), since \( V \) is continuous, \( \mu_n \to \mu \), and \( \mu_n \) and \( \mu \) are supported in the same compact set, so in fact we have established that

\[
\limsup_{n \to +\infty} \frac{1}{n^2} H_n(\mu_n) \leq I(\mu) + o_\eta(1)
\]

and letting \( \eta \to 0 \) finally gives us the \( \Gamma \)-lim sup inequality, which concludes the proof. \( \square \)

**Remark 2.19.** Again we have not really used the fact that \( g \) is a Coulombic kernel, rather we only used the fact that \( g \) is monotone radial, positive in \( B(0,1) \) and \( \iint g(x-y) \, dx \, dy < \infty \). This shows that the result still holds for all such interaction kernels.

**Remark 2.20.** To prove the \( \Gamma \)-lim inf relation, we have only used that \( V \) is l.s.c. and bounded below. To prove the \( \Gamma \)-lim sup relation, we have assumed that \( V \) is continuous for convenience. In fact the construction works for more general \( V \)'s, for example it suffices to assume that \( V \) continuous on the set where it is finite.

We next derive the consequence of the \( \Gamma \)-convergence Proposition 2.8 given by Proposition 2.6. In order to do so, we must prove the compactness of sequences with suitably bounded energy, as in Remark 2.7.

**Lemma 2.21.** Assume that \( V \) satisfies (A1)–(A2). Let \( \{(x_1, \ldots, x_n)\}_n \) be a sequence of configurations in \( \mathbb{R}^d \), and let \( \{\mu_n\}_n \) be the associated empirical measures (defined by \( \mu_n = \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i} \)). Assume \( \{\frac{1}{n} H_n(\mu_n)\}_n \) is a bounded sequence. Then the sequence \( \{\mu_n\}_n \) is tight, and as \( n \to \infty \), it converges weakly in \( \mathcal{P}(\mathbb{R}^d) \) (up to extraction of a subsequence) to some probability measure \( \mu \).

**Proof.** The proof is completely analogous to that of Lemma 2.10. First, by assumption, there exists a constant \( C_1 \) independent of \( n \) such that \( H_n \leq C_1 n^2 \), and in view of (2.37)–(2.39) we may write, for every \( M > 0 \),

\[
C_1 \geq \iint (g(x-y) \wedge M) \, d\mu_n(x) \, d\mu_n(y) \, \frac{M}{n} + \int V \, d\mu_n
\]

\[
= \iint \left[ g(x-y) \wedge M + \frac{1}{2} V(x) + \frac{1}{2} V(y) \right] \, d\mu_n(x) \, d\mu_n(y) - \frac{M}{n}. \quad (2.46)
\]

In view of (2.9), given any constant \( C_2 > 0 \), we may find \( M \) large enough and a compact set \( K \) such that

\[
\min_{(K \times K)^\circ} \left[ g(x-y) \wedge M + \frac{1}{2} V(x) + \frac{1}{2} V(y) \right] > C_2.
\]

The rest of the proof is virtually as in Lemma 2.10. \( \square \)
To conclude, we will make the assumptions on $V$ that ensure both the $\Gamma$-convergence result Proposition 2.6 and the existence result Theorem 2.1. Since we assumed for simplicity that $V$ is continuous and finite, it suffices to assume (A2) to have (A1) and (A3).

**Theorem 2.2** (Convergence of minimizers and minima of $H_n$). Assume that $V$ is continuous and satisfies (A2). Assume that for each $n$, $\{(x_1, \ldots, x_n)\}_n$ is a minimizer of $H_n$. Then,

$$\frac{1}{n} \sum_{i=1}^n \delta_{x_i} \to \mu_0 \text{ in the weak sense of probability measures} \quad (2.47)$$

where $\mu_0$ is the unique minimizer of $I$ as in Theorem 2.1 and

$$\lim_{n \to +\infty} \frac{1}{n^2} H_n(x_1, \ldots, x_n) = I(\mu_0). \quad (2.48)$$

**Proof.** Applying the $\Gamma$-limsup part of the definition of $\Gamma$-convergence, for example to $\mu_0$, ensures that $\lim \sup_{n \to +\infty} \frac{1}{n^2} \min H_n$ is bounded above (by $I(\mu_0)$), hence in particular sequences of minimizers of $H_n$ satisfy the assumptions of Lemma 2.21. It follows that, up to a subsequence, we have $\frac{1}{n} \sum_{i=1}^n \delta_{x_i} \to \mu$ for some $\mu \in \mathcal{P}(\mathbb{R}^d)$. By Propositions 2.6 and 2.8, $\mu$ must minimize $I$, hence, in view of Theorem 2.1 it must be equal to $\mu_0$. This implies that the convergence must hold along the whole sequence. We also get (2.48) from Proposition 2.6.

In the language of statistical mechanics or mean field theory, this result gives the mean-field behavior or average behavior of ground states, and the functional $I$ is called the mean-field energy functional. It tells us that points distribute themselves macroscopically according to the probability law $\mu_0$ as their number tends to $\infty$, and we have the leading order asymptotic expansion of the ground state energy

$$\min H_n \sim n^2 \min I.$$

This is not very precise in the sense that it tells us nothing about the precise patterns they follow. Understanding this is the object of the following chapters.

### 2.5 Linking the equilibrium measure with the obstacle problem

In Section 2.3 we described the characterization of the equilibrium measure minimizing $I$ via tools of potential theory. In this section, we return to this question and connect it instead to a well-studied problem in the calculus of variations called the **obstacle problem**. This connection is not often emphasized in the literature. It is however mentioned in passing in [SaTo] and used intensively in [HenMa]. It allows us to use PDE theory results, such as methods based on the maximum principle methods and regularity theory to obtain additional information on $\mu_0$. 

2.5.1 Short presentation of the obstacle problem

The obstacle problem is generally formulated over a bounded domain $\Omega \in \mathbb{R}^d$: given an $H^1(\Omega)$ function $\psi : \Omega \to \mathbb{R}$ (called the obstacle), which is nonpositive on $\partial \Omega$, find the function that achieves

$$\min \left\{ \int_\Omega |\nabla h|^2, \ h \in H^1_0(\Omega), \ h \geq \psi \right\}. \quad (2.49)$$

For general background and motivation for this problem, see e.g. [KS,Fri,Caf-Kin].

Here the space $H^1_0(\Omega)$ is the Sobolev space of trace-zero functions which is the completion of $C^1_c(\Omega)$ ($C^1$ functions with compact support in $\Omega$) under the $H^1$-Sobolev norm $||h||_{H^1} = ||h||_{L^2} + ||\nabla h||_{L^2}$. The zero trace condition $h \in H^1_0(\Omega)$ may be replaced by different boundary conditions, e.g. a translation $h \in H^1_0(\Omega) + f$, where $f$ is a given function. Note that the minimization problem (2.49) is a convex minimization problem under a convex constraint, hence it has at most one minimizer (it is not too hard to show that the minimum is achieved, hence there actually is a unique minimizer).

An admissible function for (2.49) has two options at each point: to touch the obstacle or not (and typically uses both possibilities). If $h$ is the optimizer, the set \( \{ x \in \Omega | h(x) = \psi(x) \text{ q.e.} \} \) is closed and called the coincidence set or the contact set. It is unknown (part of the problem), and its boundary is called a “free boundary.” The obstacle problem thus belongs to the class of so-called free-boundary problems, cf. [Fri].

Trying to compute the Euler-Lagrange associated to this problem by perturbing $h$ by a small function, one is led to two possibilities depending on whether $h = \psi$ or $h > \psi$. In a region where $h > \psi$, one can do infinitesimal variations of $h$ of the form $(1-t)h + tv$ with $v$, say, smooth (this still gives an admissible function, i.e. lying above the obstacle, as soon as $t$ is small enough) which shows that $\Delta h = 0$ there (since the “functional” derivative of the Dirichlet energy is the Laplacian).

In the set where $h = \psi$, only variations of the same form $(1-t)h + tv$ but with $v \geq \psi$ (equivalent to $v \geq h$ there!) and $t \geq 0$ provide admissible functions, and this only leads to an inequality $-\Delta h \geq 0$ there. These two pieces of information can be grouped in the following more compact form:

$$\text{for all } v \text{ in } H^1_0 \text{ such that } v \geq \psi \text{ q.e., } \int_\Omega \nabla h \cdot \nabla (v-h) \geq 0. \quad (2.50)$$

This relation is called a variational inequality, and it uniquely characterizes the solution to (2.49), in particular the coincidence set is completely determined as part of the solution.

In Fig. 2.2 below, we describe a few instances of solutions to one-dimensional obstacle problems, and in Fig. 2.3 to higher dimensional obstacle problems.

The regularity theory of the solutions to obstacle problems and of their coincidence sets has been developed for many years, culminating with the work of
2.5 Linking the equilibrium measure with the obstacle problem

Caffarelli (for a review see [Caff]). This sophisticated PDE theory shows, for example, that the solution $h$ is as regular as $\psi$ up to $C^{1,1}$ [Fre]. The boundary $\partial\Sigma$ of the coincidence set is $C^{1,\alpha}$ except for cusps [Caff]. These are points of $\partial\Sigma$ at which, locally, the coincidence set can fit in the region between two parallel planes separated by an arbitrarily small distance (the smallness of the neighborhood de-
The leading order behavior of the Coulomb gas

Fig. 2.4 gives examples of coincidence sets, one regular one, and one with cusps.

Moreover, if \( \psi \) is \( C^{1,1} \), since \( \nabla h \) is continuous, the graph of \( h \) must leave the coincidence set tangentially. This formally leads to the following system of equations, where \( \omega \) denotes the coincidence set:

\[
\begin{align*}
-\Delta h &= 0 \text{ in } \Omega \setminus \omega \\
h &= \psi \quad \text{in } \omega \\
\frac{\partial h}{\partial \nu} &= \frac{\partial \psi}{\partial \nu} \quad \text{on } \partial \omega \\
h &= 0 \quad \text{on } \Omega.
\end{align*}
\]

This relation cannot be made rigorous in all cases, because \( \omega \) is not an open domain, however it gives the right intuition and is correct when \( \omega \) is nice enough. Note that on the boundary of \( \Omega \setminus \omega \) we must have a Dirichlet condition \( h = \psi \), together with a Neumann condition \( \frac{\partial h}{\partial \nu} = \frac{\partial \psi}{\partial \nu} \). These two boundary conditions make what is called an overdeternimented problem and this overdetermination explains why there is only one possible coincidence set.

### 2.5.2 Connection between the two problems

The problem we examined, that of the minimization of \( I \), is phrased in the whole space, and not in a bounded domain. While the minimization problem (2.49) may not have a meaning over all \( \mathbb{R}^d \) (because the integral might not converge), the corresponding variational inequality (2.50) can still be given a meaning over \( \mathbb{R}^d \) as follows: given \( \psi \in H^1_{\text{loc}}(\mathbb{R}^d) \) solve for

\[
\forall v \in \mathcal{K}, \quad \int_{\mathbb{R}^d} \nabla h \cdot \nabla(v - h) \geq 0 \tag{2.51}
\]

where

\[
\mathcal{K} = \left\{ H^1_{\text{loc}}(\mathbb{R}^d) \text{ such that } v - h \text{ has bounded support and } v \geq \psi \text{ q.e.} \right\}.
\]
Solving this is in fact equivalent to the statement that for every $R > 0$, $h$ is the unique solution to

$$
\min \left\{ \int_{B_R} |\nabla v|^2, v \in H^1(B_R), v - h \in H^1_0(B_R), v \geq \psi \text{ in } B_R \right\},
$$

which replaces (2.49). The problem (2.51) is easily seen to have a unique solution: if there are two solutions $h_1$ and $h_2$ it suffices to apply (2.51) for $h_1$, with $h_2$ as a test-function, and then reverse the roles of the two and add the two relations to obtain $h_1 = h_2$.

Let us now compare the two problems side by side:

**Equilibrium measure.** $\mu_0$ is characterized by the relations

$$
\begin{cases}
    h^\mu_0 + \frac{V}{2} \geq c & \text{quasi everywhere} \\
    h^\mu_0 + \frac{V}{2} = c & \text{q.e. in the support of } \mu_0.
\end{cases}
$$

(2.52)

**Obstacle problem.**

$$
\begin{cases}
    h \geq \psi & \text{q.e.} \\
    h = \psi & \text{q.e. in the coincidence set}
\end{cases}
$$

(2.53)

It is then not surprising to expect a correspondence between the two settings, once one chooses the obstacle to be $\psi = c - \frac{V}{2}$.

**Proposition 2.22** (Equivalence between the minimization of $I$ and the obstacle problem). Assume $d \geq 2$, $V$ is continuous and satisfies (A2). If $\mu_0$ is the equilibrium measure associated to the potential $V$ as in Theorem 2.1 then its potential $h^\mu_0$, as defined in (2.16), is the unique solution to the obstacle problem with obstacle $\psi = c - \frac{V}{2}$ in the sense of (2.51). If in addition $V \in C^{1,1}$ then $\mu_0 = (\frac{1}{2} \Delta V)_{1,\omega}$.

Note that the converse might not be true, because a solution of the obstacle problem can fail to provide a probability measure, however it does in general when shifting $c$ appropriately.

When one works on a bounded domain, this result can be obtained by observing that the problem of minimizing $I$ and that of minimizing (2.49) are essentially convex duals of each other (see [Bre,BrSer]). When working in an infinite domain, the correspondence is probably folklore and could also be worked out by convex duality, but we were not able to find it completely written in the literature, except for [HenMa] who follow a slightly different formulation. Here, for the proof, we follow the approach of [ASZ] where the result is established in dimension 2 for the particular case of $V$ quadratic (but with more general constraints), the adaptation to any dimension and to general $V$’s is not difficult.

**Proof. Step 1.** We show that $\nabla h^\mu_0$ is in $L^2_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d)$. It is a consequence of the fact that $I(\mu_0) < \infty$ hence, in view of the assumptions on $V$, $\int g(x -$
which smooth measures, and combining Corollary 5.10 and Theorem 7.9 in [LiLo].

In the case $d = 2$, we need to consider a reference probability measure $\bar{\mu}$ for which $h^0$ is $C^1_{loc}(\mathbb{R}^2)$. It suffices to consider for example $\bar{\mu} = \frac{\mathbb{1}}{2}1_{B_r}$, the circle law, for which $h^0$ is radial and can be computed explicitly. Then, let us consider $\rho = \mu_0 - \bar{\mu}$. Using the fact that $\int d\rho = 0$, $\rho$ is compactly supported, and $\int g(x - y) d\rho(x) d\rho(y) < \infty$ holds for both $\mu = \mu_0$ and $\mu = \bar{\mu}$, we have the following statement

$$\int_{\mathbb{R}^2 \times \mathbb{R}^2} -\log |x - y| d\rho(x) d\rho(y) = \frac{1}{2\pi} \int_{\mathbb{R}^2} |\nabla h^\rho(x)|^2 dx,$$

(2.54)

where $h^\rho(x) = \int g(x - y) d\rho(y)$. Indeed, in the proof of [SaTo, Lemma 1.8] it is shown that

$$\int_{\mathbb{R}^2} \int_{\mathbb{R}^2} -\log |x - y| d\rho(x) d\rho(y) = \frac{1}{2\pi} \int_{\mathbb{R}^2} \left( \int_{\mathbb{R}^2} \frac{1}{|x - y|} d\rho(y) \right)^2 dx,$$

and (2.54) follows, since $\int_{\mathbb{R}^2} \frac{(x - y) d\rho(y)}{|x - y|^2} = -\nabla h^\rho(x)$ in the distributional sense. This shows that $\nabla h^\rho \in L^2(\mathbb{R}^2)$ and thus, since $h^\rho$ is $C^1_{loc}(\mathbb{R}^2)$, we deduce that $\nabla h^\mu = \nabla h^\rho + \nabla h^\bar{\mu}$ is also in $L^2_{loc}(\mathbb{R}^2, \mathbb{R}^2)$, as desired.

**Step 2.** Let $v$ be admissible in (2.51), i.e. belong to $K$, and set $\varphi = v - h^\mu_0$. If $\varphi$ is smooth and compactly supported, then

$$\int_{\mathbb{R}^d} \nabla h^\mu_0 \cdot \nabla (v - h^\mu_0) = c_d \int_{\mathbb{R}^d} \varphi d\mu_0 \geq 0. \quad (2.55)$$

Indeed, by (2.15), we know that $h^\mu_0 = \psi$ q.e. in the support of $\mu_0$ and by assumption $v \geq \psi$ q.e. in $\mathbb{R}^d$. Hence $\varphi$ is q.e. nonnegative on the support of $\mu_0$ and the inequality (2.55) follows, since $\mu_0$ does not charge sets of zero capacity. To obtain (2.55) for any $v \in K$, it suffices to show that the subset of $K$ consisting of $v$'s for which $v - h^\mu_0$ is smooth and compactly supported is dense in $K$ for the topology of $H^1$. Fix some $v$ in the admissible set and $R > 1$ such that $v - h^\mu_0$ is supported in $B_R/2$. Let $\eta_\varepsilon$ be a standard mollifier and $\chi_R$ a smooth function supported in $B_{2R}$ with $0 \leq \chi_R \leq 1$ and $\chi_R \equiv 1$ in $B_R$. One may check that

$$v_{\varepsilon, \delta} = h^\mu_0 + (v - h^\mu_0) * \eta_\varepsilon \chi_R$$

satisfies that $v_{\varepsilon, \delta} - h^\mu_0$ is smooth and approximates $v$ arbitrarily well in $H^1$ when $\delta$ is small enough, and is $\geq \psi$ when $\varepsilon$ is chosen small enough relative to $\delta$. This concludes the proof of (2.55).

**Step 3.** We prove the statements about $\mu_0$. First, since the coincidence set $\omega$ is closed, its complement is open, and the function $h^\mu_0$ is harmonic on that set. One can note also that in view of (2.15) and the definition of the coincidence set $\omega$, the support of $\mu_0$ is included in $\omega$ up to a set of zero capacity.
If we assume that \( V \in C^{1,1}_{\text{loc}} \), then by Frehse’s regularity theorem mentioned above, it follows that \( h^{\mu_0} \) is also \( C^{1,1}_{\text{loc}} \). In particular \( h^{\mu_0} \) is continuous, and so is \( V \), so the relations \((2.15)\) hold pointwise and not only q.e. This means that we have
\[
h^{\mu_0} + \frac{1}{2} V = c \quad \text{on } \omega
\] (2.56)
and \( \text{Supp}(\mu_0) \subset \omega \). Also \( C^{1,1}_{\text{loc}} = W^{2,\infty}_{\text{loc}} \) hence \( \Delta h^{\mu_0} \) and \( \Delta V \) both make sense as \( L^\infty_{\text{loc}} \) functions, and it suffices to determine \( \mu_0 \) up to sets of measure 0. We already know that \( \mu_0 = 0 \) in the complement of \( \omega \) since \( h^{\mu_0} \) is harmonic there, and it suffices to determine it in \( \overset{\circ}{\omega} \). But taking the Laplacian on both sides of \((2.56)\), since \( \mu_0 = -\frac{1}{c_d} \Delta h^{\mu_0} \), one finds
\[
\mu_0 = \frac{1}{2c_d} \Delta V \quad \text{in } \overset{\circ}{\omega},
\]
and the results follows.

\[\square\]

By definition of \( \zeta \) \((2.36)\), we have that
\[
\{ x \in \mathbb{R}^d | \zeta(x) = 0 \} = \omega.
\] (2.57)

Since \( \mu_0 \) is a compactly supported probability measure, we have that \( h^{\mu_0} = \int g(x-y) \, d\mu_0(y) \) asymptotically behaves like \( g(x) \) as \( |x| \to \infty \). Since \( h^{\mu_0} + \frac{1}{2} V = c \) q.e. in \( \omega \) and since \((A2)\) holds, it follows that \( \omega \) must be a bounded, hence compact, set.

We have seen that \( \omega \) contains, but is not always equal to, the support of \( \mu_0 \). The latter is called the droplet in \( \text{[HenMa]} \), where similar results to this proposition are established. There, it is also discussed how \( \text{Supp}(\mu_0) \) differs from \( \omega \) (they are equal except at “shallow points”, cf. definition there).

**Remark 2.23** (Note on dimension one). For \( d = 1 \) and \( g = -\log |\cdot| \), as seen before \( h^{\mu_0} \) solves
\[
- \Delta^{1/2} h^{\mu_0} = c_1 \mu_0,
\]
and the equivalent of the obstacle problem is instead a fractional obstacle problem for which a good theory also exists \( \text{[CSS]} \). One could also write the analogue of Proposition \( 2.22 \).

We have seen how the correspondence between the minimization of \( I \) and the obstacle problem thus allows, via the regularity theory of the obstacle problem, to identify the equilibrium measure in terms of \( V \) when the former is regular enough. The known techniques on the obstacle problem \( \text{[Caff]} \) also allow for example to analyze the rate at which the solution leaves the obstacle (they say it is subquadratic), which gives us information on the size of the function \( \zeta \), defined in \((2.36)\).
2.6 Large deviations for the Coulomb gas with temperature

At this point, we know the $\Gamma$-convergence of $\frac{1}{n^2} H_n$ and its consequence, Theorem 2.24 for ground states of the Coulomb gas. In this section, we turn for the first time to states with temperature and derive rather easy consequences of the previous sections on the Gibbs measure, which we recall is defined by

$$d\mathbb{P}_{n,\beta}(x_1, \ldots, x_n) := \frac{1}{Z_{n,\beta}} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n$$

(2.58)

with

$$Z_{n,\beta} = \int e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n.$$  

(2.59)

Pushing $\mathbb{P}_{n,\beta}$ forward by the map $(x_1, \ldots, x_n) \mapsto \frac{1}{n} \sum_{i=1}^n \delta_{x_i}$, we may view it as a probability measure on $\mathcal{P}(\mathbb{R}^d)$, called the Gibbs measure at (inverse) temperature $\beta$. Note that considering $\beta$ that depends on $n$ can correspond to other temperature regimes (very high or very low temperatures) and is also interesting.

Formally, taking $\beta = +\infty$ above reduces to the study to the minimizers of $H_n$, whose behavior when $n \to +\infty$ we already established (weak convergence of the empirical measure to the equilibrium measure $\mu_0$). For $\beta < +\infty$ we will see that the behavior of a “typical” configuration $(x_1, \ldots, x_n)$ under the measure $\mathbb{P}_{n,\beta}$ is not very different, and we can even characterize the probability of observing a “non-typical” configuration. The sense given to “typical” and “non-typical” will be that of the theory of large deviations, which we first briefly introduce. For more reference, one can see the textbooks [DenH,DS,DZ].

**Definition 2.24** (Rate function). Let $X$ be a metric space (or a topological space). A rate function is a l.s.c. function $I: X \to [0, +\infty]$, it is called a “good rate function” if its sub-level sets $\{x, I(x) \leq \alpha\}$ are compact (see Remark 2.2).

**Definition 2.25** (Large deviations). Let $\{P_n\}_n$ be a sequence of Borel probability measures on $X$ and $\{a_n\}$ a sequence of positive real numbers diverging to $+\infty$. Let also $I$ be a (good) rate function on $X$. The sequence $\{P_n\}_n$ is said to satisfy a large deviation principle (LDP) at speed $a_n$ with (good) rate function $I$ if for every Borel set $E \subset X$ the following inequalities hold:

$$-\inf_E I \leq \liminf_{n \to +\infty} \frac{1}{a_n} \log P_n(E) \leq \limsup_{n \to +\infty} \frac{1}{a_n} \log P_n(E) \leq -\inf_{\overline{E}} I$$

(2.60)

where $\overline{E}$ (resp. $\overline{E}$) denotes the interior (resp. the closure) of $E$ for the topology of $X$.

Formally, it means that $P_n(E)$ should behave roughly like $e^{-a_n \inf_E I}$. The rate function $I$ is the rate of exponential decay of the probability of rare events, and the events with larger probability are the ones on which $I$ is smaller.
Remark 2.26. At first sight, Definition 2.25 looks very close to the $\Gamma$-convergence

$$\frac{\log p_n}{a_n} \to -I$$

where $p_n$ is the density of the measure $P_n$. However, in general there is no equivalence between the two concepts. For example, in order to estimate the quantity

$$\log P_n(E) = \log \int_E p_n(x) dx$$

(2.61)

it is not sufficient to know the asymptotics of $p_n$, one really also needs to know the size of the volume element $\int_E dx$, which plays a large role in large deviations and usually comes up as an entropy term. There are however some rigorous connections between $\Gamma$-convergence and LDP (see [Mar]).

We will need an additional assumption on $V$:

(A4) There exists $\alpha > 0$ such that

$$\int_{\mathbb{R}^d} e^{-\alpha V(x)} dx < +\infty.$$  (2.62)

We will also keep the other assumptions that $V$ is continuous and (A2) holds, which ensure the existence of the equilibrium measure $\mu_0$, and the $\Gamma$-convergence of $\frac{H_n}{a_n}$ to $I$. In dimension $d = 2$, the growth assumption (A2) $\frac{V}{2} - \log \to +\infty$ ensures that the condition (A4) is also satisfied, however in dimension $d \geq 3$ we need to assume (2.62), which is a slight strengthening of (A2), in order to avoid very slow divergence of $V$ such as $V(x) \sim \log \log x$ at infinity. Note in particular that (A4) ensures that the integral in (2.59) is convergent, hence $Z_{n,\beta}$ well-defined, as soon as $n$ is large enough.

We may now state the LDP for the Gibbs measure associated to the Coulomb gas Hamiltonian. This result is due to [PoHi] (in dimension 2), [BG] (in dimension 1) and [BZ] (in dimension 2) for the particular case of a quadratic potential (and $\beta = 2$), see also [Ber] for results in a more general (still determinantal) setting of multidimensional complex manifolds [CGZ] recently treated more general singular $g$’s and $V$’s. We present here the proof for the Coulomb gas in any dimension and general potential, which is not more difficult.

Theorem 2.3 (Large deviations principle for the Coulomb gas at speed $n^2$).

Assume $V$ is continuous and satisfies (A2) and (A4). For any $\beta > 0$, the sequence $\{P_{n,\beta}\}$ of probability measures on $\mathcal{P}(\mathbb{R}^d)$ satisfies a large deviations principle at speed $n^2$ with good rate function $\frac{\beta}{2} I$ where $I = I - \min_{\mathcal{P}(\mathbb{R}^d)} I = I - I(\mu_0)$. Moreover

$$\lim_{n \to +\infty} \frac{1}{n^2} \log Z_{n,\beta} = -\frac{\beta}{2} I(\mu_0) = -\frac{\beta}{2} \min_{\mathcal{P}(\mathbb{R}^d)} I.$$  (2.63)
Here of course, the underlying topology is still that of weak convergence on $\mathcal{P}(\mathbb{R}^d)$.

The heuristic reading of the LDP is that

$$
P_{n,\beta}(E) \approx e^{-\frac{1}{2n^2} \min_E I - \min_I}.
$$

(2.64)

As a consequence, the only likely configurations of points (under $P_{n,\beta}$) are those for which the empirical measures $\mu_n = \frac{1}{n} \sum_{i=1}^{n} \delta_{x_i}$ converge to $\mu = \mu_0$, for otherwise $I(\mu) > I(\mu_0)$ by uniqueness of the minimizer of $I$, and the probability decreases exponentially fast according to (2.64). Thus, $\mu_0$ is not only the limiting distribution of minimizers of $H_n$, but also the limiting distribution for all “typical” (or likely) configurations, at finite temperature. Moreover, we can estimate the probability under $P_{n,\beta}$ of the “non-typical” configurations and see that it has exponential decay at speed $n^2$. Recall that the cases of the classic random matrix ensembles GOE, GUE and Ginibre correspond respectively to $d = 1$, $\beta = 1, 2$ and $V(x) = x^2$, and $d = 2$, $\beta = 2$, and $V(x) = |x|^2$. The corresponding equilibrium measures were given in Example 2 above. As a consequence of Theorem 2.3 we have a proof that the distribution of eigenvalues (more precisely the spectral or empirical measure) has to follow Wigner’s semi-circle law $\mu_0 = \frac{1}{\pi} \sqrt{4 - x^2} 1_{|x| < 2}$ for the GUE and GOE, and the circle law $\mu_0 = \frac{1}{\pi} 1_{B_1}$ for the Ginibre ensemble, in the sense of the LDP (which is in fact stronger than just establishing these laws). These are the cases originally treated in [PeHi,BZ,BG].

In addition, knowing the partition function $Z_{n,\beta}$ is important because it gives access to many physical quantities associated to the system (for e.g. by differentiating $Z_{n,\beta}$ with respect to $\beta$ yields the average energy, etc), see statistical mechanics textbooks such as [Huan]. In particular $-\frac{2}{\beta} \log Z_{n,\beta}$ in our context is physically the free energy of the system, and the existence of a limit for $\log Z_{n,\beta}$ properly normalized is called the existence of thermodynamic limit. For the one-dimensional log gas, the value of $Z_{n,\beta}$ is known explicitly for all $\beta > 0$ when $V(x) = x^2$ via the exact computation of the integral in (2.59), which uses so-called Selberg integrals (see e.g. [Me]). For more general $V$’s an expansion in $n$ to any order is also known [BoGui]. In dimension 2 however, no equivalent of the Selberg integral exists and the exact value of $Z_{n,\beta}$ is only known for the Ginibre case $\beta = 2$ and $V(x) = |x|^2$ (there are a few other exceptions). In dimension $d \geq 3$, we know of no such explicit computation or expansion.

**Proof of the theorem.** The intuition behind the LDP might be that since the sequence $\frac{1}{n^2} H_n$ $\Gamma$-converges to $I$, we should have

$$
P_{n,\beta} \approx \frac{1}{Z_n} e^{-\frac{\beta}{2} n^2 I}
$$

however such an approach is too naive to work directly, for the reasons explained in Remark 2.26. We will use the $\Gamma$-convergence result in a more precise way, also estimating the size of the appropriate sets in configuration space.
Step 1. We first prove the large deviations upper bound, that is:

\[
\limsup_{n \to +\infty} \frac{\log \mathbb{P}_{n,\beta}(E)}{n^2} \leq -\inf_E \tilde{I}
\]  

(2.65)

up to an estimate on \(\log Z_{n,\beta}\), then we will turn to the proof of the lower bound and get in passing the missing estimate on \(\log Z_{n,\beta}\).

Let us define \(\tilde{H}_n(x_1, \ldots, x_n) = H_n(x_1, \ldots, x_n) - \sqrt{n} \sum_{i=1}^n V(x_i)\). This amounts to changing \(V\) to \((1 - \frac{1}{\sqrt{n}})V\) in the definition of \(H_n\). Of course, \(\tilde{H}_n\) still \(\Gamma\)-converges to \(I\), by the same proof as Proposition 2.8. We want to show that

\[
\liminf_{n \to +\infty} \frac{\hat{H}_n}{n^2} \geq \inf_E \hat{I}.
\]  

(2.66)

We may assume that the left-hand side is finite, otherwise there is nothing to prove. Upon passing to a subsequence, suppose that \(\mu_n = \frac{1}{n} \sum_{i=1}^n \delta_{x_i}\) is a minimizer (or almost minimizer) of \(\hat{H}_n\) on \(E\), more precisely satisfies

\[
\lim_{n \to +\infty} \frac{\hat{H}_n(\mu_n)}{n^2} = \liminf_{n \to +\infty} \frac{\hat{H}_n}{n^2}.
\]

By applying Lemma 2.21 (applied to \(\hat{H}_n\) instead of \(H_n\) but this induces no change), the sequence \(\{\mu_n\}_n\) is tight and has a subsequence which converges to some \(\mu\) in the sense of weak convergence in \(\mathcal{P}(\mathbb{R}^d)\). Then by the \(\Gamma\)-convergence of \(\frac{\hat{H}_n}{n^2}\) to \(I\), it follows that

\[
\liminf_{n \to +\infty} \frac{\hat{H}_n(\mu_n)}{n^2} = I(\mu) \geq \inf_E I
\]

and (2.66) is proven. In particular, (2.66) implies that \(\frac{1}{n^2} \hat{H}_n \geq \inf_E I + o(1)\) on \(E\), the \(o(1)\) being uniform on \(E\). Inserting this inequality into the definition of \(\mathbb{P}_{n,\beta}\), one gets

\[
\mathbb{P}_{n,\beta}(E) \leq \frac{1}{Z_{n,\beta}} \int_{\mathbb{R}^d} e^{-\frac{\beta}{2} n^2 \inf_{\mathbb{R}^d} I + o(1)} e^{-\frac{\beta}{2} \sqrt{n} \sum_{i=1}^n V(x_i)} dx_1 \ldots dx_n
\]

\[
= \frac{1}{Z_{n,\beta}} e^{-\frac{\beta}{2} n^2 \left(\inf_{\mathbb{R}^d} I + o(1)\right)} \left(\int_{\mathbb{R}^d} e^{-\frac{\beta}{2} \sqrt{n} V(x)} dx\right)^n.
\]  

(2.67)

The last integral in the right-hand side is bounded by a constant for \(n\) large enough by the assumption (A4). Hence, taking the logarithm of both sides, we find

\[
\log \mathbb{P}_{n,\beta}(E) \leq -\log Z_{n,\beta} - \frac{\beta}{2} n^2 \inf_E I + o(n^2) + O(n) \text{ as } n \to +\infty,
\]  

(2.68)

and the \(\limsup\) of (2.68) when \(n \to +\infty\) gives, for each \(E \subset \mathcal{P}(\mathbb{R}^d)\)

\[
\limsup_{n \to +\infty} \frac{1}{n} \log \mathbb{P}_{n,\beta}(E) \leq \limsup_{n \to +\infty} \left( -\frac{1}{n^2} \log Z_{n,\beta} \right) - \frac{\beta}{2} \inf_E I.
\]  

(2.69)
This is not exactly the large deviations upper bound relation, because we cannot yet bound the term $-\frac{1}{n^2} \log Z_{n,\beta}$. However, by taking $E$ to be the whole space $\mathcal{P}(\mathbb{R}^d)$ in (2.68), we already get that
\[
\frac{1}{n^2} \log Z_{n,\beta} \leq -\frac{\beta}{2} I(\mu_0) + o(1) \quad \text{as} \quad n \to +\infty. \tag{2.70}
\]

**Step 2.** We prove the large deviations lower bound. Let $\mu$ be in the interior $\overset{\circ}{E}$. By the $\Gamma$-lim sup result of Proposition 2.8, there exists a sequence of $n$-tuples $(x_1, \ldots, x_n)$ such that the empirical measures $\frac{1}{n} \sum_{i=1}^{n} \delta_{x_i}$ converges weakly to $\mu$ and
\[
\limsup_{n \to +\infty} \frac{H_n}{n^2}(x_1, \ldots, x_n) \leq I(\mu). \tag{2.71}
\]
Moreover, in the construction for the $\Gamma$-limsup made in Step 2 of the proof of Proposition 2.8, the balls $B(x_i, \frac{\lambda}{n^{1/2}})$ were disjoint. Consequently, if $(y_1, \ldots, y_n)$ is such that for each $i = 1 \ldots n$, the point $y_i$ is in the (Euclidean) ball $B(x_i, \frac{\lambda}{n^{1/2}})$, then the same will hold with $2\lambda$ instead of $4\lambda$, and one can check that the proof carries through in the same way, yielding that the empirical measure $\frac{1}{n} \sum_{i=1}^{n} \delta_{y_i}$ converges weakly to $\mu$ and we have
\[
\limsup_{n \to +\infty} \frac{H_n}{n^2}(y_1, \ldots, y_n) \leq I(\mu). \tag{2.72}
\]

This is helpful because it shows that there are configurations whose Hamiltonian is not too large (the upper bound (2.72) holds) and, as shown below, that these configurations occupy enough volume in phase space to contribute significantly to the partition function $Z_{n,\beta}$. Denote by $V$ the set
\[
V := \bigcup_{\sigma \in S_n} \prod_{i=1}^{n} B(x_{\sigma(i)}, \frac{\lambda}{n^{1/2}}) \tag{2.73}
\]
where $S_n$ is the set of all permutations of $\{1, \cdots, n\}$. Since the Hamiltonian $H_n$ is symmetric, it is invariant under the action of permutation of points in phase space. Thus, in view of (2.72), we have
\[
\limsup_{n \to +\infty} \max_V H_n \leq I(\mu). \tag{2.74}
\]
Moreover, for $n$ large enough the measures $\frac{1}{n} \sum_{i=1}^{n} \delta_{y_i}$ (with $(y_1, \ldots, y_n) \in V$) are in $E$, because $\frac{1}{n} \sum_{i=1}^{n} \delta_{y_i}$ converges weakly to $\mu \in \overset{\circ}{E}$. Therefore, we may write
\[
\mathbb{P}_{n,\beta}(E) \geq \frac{1}{Z_{n,\beta}} \int_{V} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n
\]
\[
= \frac{|V|}{Z_{n,\beta}} e^{-\frac{\beta}{2} (n^2 I(\mu) + o(n^2))} \quad \text{as} \quad n \to +\infty. \tag{2.75}
\]
2.6 Large deviations for the Coulomb gas with temperature

We need to estimate the volume $|\mathcal{V}|$ of $\mathcal{V}$, and it is easy to see that (for a certain constant $C$ depending only on the dimension)

$$|\mathcal{V}| = n! \left( \frac{C \lambda^d}{n} \right)^n$$

hence

$$\log |\mathcal{V}| = \log n! - n \log n + O(n) = O(n), \quad (2.76)$$

where we used that, by Stirling’s formula, $\log n! = n \log n + O(n)$. Taking the logarithm of (2.75) and inserting (2.76), one gets

$$\frac{1}{n^2} \log \mathbb{P}_{n,\beta}(E) \geq - \frac{1}{n^2} \log Z_{n,\beta} - \frac{\beta}{2} I(\mu) + o(1). \quad (2.77)$$

Since this is true for any $\mu \in \hat{E}$, taking the supremum with respect to $\mu \in \hat{E}$ and the $\liminf_{n \to +\infty}$ of (2.77) gives, for each $E \subset \mathcal{P}(\mathbb{R}^d)$:

$$\liminf_{n \to +\infty} \frac{1}{n^2} \log \mathbb{P}_{n,\beta}(E) \geq \liminf_{n \to +\infty} \left( - \frac{1}{n^2} \log Z_{n,\beta} \right) - \frac{\beta}{2} \inf_E I. \quad (2.78)$$

Moreover, choosing $E = \mathcal{P}(\mathbb{R}^d)$ and $\mu = \mu_0$ in (2.77), we obtain

$$\frac{\log Z_{n,\beta}}{n^2} \geq - \frac{\beta}{2} I(\mu_0) + o(1). \quad (2.79)$$

Combining the two inequalities (2.70) and (2.79), we get the second part of the theorem, that is the thermodynamic limit at speed $n^2$:

$$\lim_{n \to +\infty} \frac{\log Z_{n,\beta}}{n^2} = - \frac{\beta}{2} I(\mu_0) = - \frac{\beta}{2} \min_{\mathcal{P}(\mathbb{R}^d)} I. \quad (2.80)$$

Finally, inserting (2.80) into (2.69) and (2.78) completes the proof of Theorem 2.3. \qed
3 The next order behavior: splitting the Hamiltonian

In Chapter 2, we examined the leading order behavior of the Hamiltonian $H_n$ of the Coulomb gas, which can be summarized by:

- The minimal energy $\min H_n$ behaves like $n^2 \min I$, where $I$ is the mean-field limit energy, defined on the set of probability measures of $\mathbb{R}^d$.
- If each $(x_1, \ldots, x_n)$ minimizes $H_n$, the empirical measures $\frac{1}{n} \sum_{i=1}^n \delta_{x_i}$ converge weakly to the unique minimizer $\mu_0$ of $I$, also known as Frostman’s equilibrium measure, which can be characterized via an obstacle problem.
- This behavior also holds when $\beta < +\infty$, except with a very small probability determined by a large deviation principle.

The following questions thus arise naturally:

1. What lies beyond the term $n^2 I(\mu_0)$ in the asymptotic expansion of $\min H_n$ as $n \to +\infty$ and in the expansion of the partition function $\log Z_{n,\beta} = \frac{\beta}{2} n^2 I(\mu_0) + o(n^2)$? Is the next term of order $n$?

2. What is the optimal microscopic distribution of the points?

To study these questions, we wish to zoom or blow-up the configurations by the factor $n^{1/d}$ (the inverse of the typical distance between two points), so that the points are well-separated (typically with distance $O(1)$), and find a way of expanding the Hamiltonian to next order. This will occupy the next chapters.

Henceforth, we will keep the same notation as in Chapter 2 and in this chapter and the following, we will make the following assumptions: $V$ is such that the unique Frostman equilibrium measure $\mu_0$ exists (for example, as we saw it suffices to require that $V$ is continuous and satisfies $(A3)$), and $\mu_0$ is absolutely continuous with respect to the Lebesgue measure, with a density $m_0(x)$ which is bounded above, or in $L^\infty(\mathbb{R}^d)$. By abuse of notation, we will often write $\mu_0(x)$ instead of $m_0(x)$.

3.1 Splitting the Hamiltonian

The splitting consists in an exact formula that separates the leading $(n^2)$ order term in $H_n$ from next order terms.

The starting point is the following: given a configuration of points $(x_1, \ldots, x_n) \in (\mathbb{R}^d)^n$, let us set $\nu_n = \sum_{i=1}^n \delta_{x_i}$ (which is not a probability measure anymore, but merely a purely atomic Radon measure). Since we expect $\frac{1}{n} \nu_n$ to converge to $\mu_0$, let us expand $\nu_n$ as

$$\nu_n = n \mu_0 + (\nu_n - n \mu_0).$$

(3.1)
3.1 Splitting the Hamiltonian

The first term in the right-hand side gives the leading order, and the second one describes the fluctuation of \( \nu_n \) around it. Note that in contrast to the equilibrium measure \( \mu_0 \) assumed to be a nice measure with a bounded density, the fluctuation \( \nu_n - n\mu_0 \) is still singular, with an atom at each point of the configuration.

Inserting the splitting (3.1) into the definition of \( H_n \), one finds

\[
H_n(x_1, \ldots, x_n) = \sum_{i \neq j} g(x_i - x_j) + n \sum_{i=1}^{n} V(x_i)
\]

\[
= \int_{\triangle^c} g(x - y) d\nu_n(x) d\nu_n(y) + n \int V d\nu_n
\]

\[
= n^2 \int_{\triangle^c} g(x - y) d\mu_0(x) d\mu_0(y) + n^2 \int V d\mu_0
\]

\[
+ 2n \int_{\triangle^c} g(x - y) d\mu_0(x) d(\nu_n - n\mu_0)(y) + n \int V d(\nu_n - n\mu_0)
\]

\[
+ \int_{\triangle^c} g(x - y) d(\nu_n - n\mu_0)(x) d(\nu_n - n\mu_0)(y).
\]

(3.2)

We now recall that \( \zeta \) was defined in (2.36) by

\[
\zeta = h\mu_0 + \frac{V}{2} - c = \int g(x - y) d\mu_0(y) + \frac{V}{2} - c
\]

(3.3)

and that \( \zeta = 0 \) in \( \Sigma \) (with the assumptions we made, one can check that \( \zeta \) is continuous, so the q.e. relation can be upgraded to everywhere).

With the help of this we may rewrite the medium line in the right-hand side of (3.2) as

\[
2n \int_{\triangle^c} g(x - y) d\mu_0(x) d(\nu_n - n\mu_0)(y) + n \int V d(\nu_n - n\mu_0)
\]

\[
= 2n \int (h\mu_0 + \frac{V}{2}) d(\nu_n - n\mu_0) = 2n \int (\zeta + c) d(\nu_n - n\mu_0)
\]

\[
= 2n \int \zeta d\nu_n - 2n^2 \int \zeta d\mu_0 + 2nc \int d(\nu_n - n\mu_0) = 2n \int \zeta d\nu_n.
\]

The last equality is due to the facts that \( \zeta \equiv 0 \) on the support of \( \mu_0 \) and that \( \nu_n \) and \( n\mu_0 \) have the same mass \( n \). We also have to notice that since \( \mu_0 \) has a \( L^\infty \) density with respect to the Lebesgue measure, it does not charge the diagonal \( \Delta \) (whose Lebesgue measure is zero) and we can include it back in the domain of integration. By that same argument, one may recognize in the first line of the right-hand side of (3.2), the quantity \( n^2 I(\mu_0) \), cf. (2.7).
We may thus rewrite (3.2) as

\[ H_n(x_1, \ldots, x_n) = n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) + \int_{\Sigma^c} g(x-y)d(\nu_n - n\mu_0)(x)d(\nu_n - n\mu_0)(y). \] (3.4)

Note that this is an exact relation, valid for any configuration of points. The first term in the right-hand side gives the leading order, i.e. the energy of the equilibrium measure. In the second term, \( \zeta \) plays the role of an effective confining potential, which is active only outside of \( \Sigma \) (recall \( \zeta \geq 0 \), and \( \zeta = 0 \) in \( \Sigma \)). The last term in the right-hand side is the most interesting, it measures the discrepancy between the diffuse equilibrium measure \( \mu_0 \) and the discrete empirical measure \( \frac{1}{n} \nu_n \). It is an electrostatic (Coulomb) interaction between a “negatively charged background” \( -n\mu_0 \) and the \( n \) positive discrete charges at the points \( x_1, \ldots, x_n \). In the sequel, we will express this energy term in another fashion, and show that it is indeed a lower-order term.

To go further, we need to introduce \( h_n \), the potential generated by the distribution of charges \( \nu_n - n\mu_0 \), defined by

\[ h_n := g* (\nu_n - n\mu_0) = \int g(x-y)d(\nu_n - n\mu_0)(y). \] (3.5)

In dimension \( d \geq 2 \) this is equivalent to

\[ h_n = -c_d \Delta^{-1}(\nu_n - n\mu_0), \] (3.6)

and in dimension 1 to

\[ h_n = -c_1 \Delta^{-1/2}(\nu_n - n\mu_0). \] (3.7)

Note that \( h_n \) decays at infinity, because the charge distribution \( \nu_n - n\mu_0 \) is compactly supported and has zero total charge, hence, when seen from infinity behaves like a dipole. More precisely, \( h_n \) decays like \( \nabla g \) at infinity, that is \( O(\frac{1}{r}) \) and its gradient \( \nabla h_n \) decays like the second derivative \( D^2 g \), that is \( O(\frac{1}{r}) \) (in dimension 1, like \( 1/r \) and \( 1/r^2 \)). Formally, using Green’s formula (or Stokes’ theorem) and the definitions, one would like to say that, at least in dimension \( d \geq 2 \),

\[ \int_{\Sigma^c} g(x-y)d(\nu_n - n\mu_0)(x)d(\nu_n - n\mu_0)(y) = \int h_n d(\nu_n - n\mu_0) \]
\[ = \int h_n (-c_d \Delta h_n) \approx c_d \int |\nabla h_n|^2 \] (3.8)

This is the place where we really use for the first time in a crucial manner the Coulombic nature of the interaction kernel \( g \). Such a computation allows to replace the sum of pairwise interactions of all the charges and “background” by an integral
3.2 A precursor to the renormalized energy

The goal of this section is to give a rigorous meaning to the formal computation (3.8). We restrict for now to the case of \( d \geq 2 \). The case \( d = 1 \) will be dealt with in the next section.

We next define the right quantity to make sense of (3.8).

Definition 3.1 (Precursor to the renormalized energy). Let \( \Omega \) be an open set in \( \mathbb{R}^d \) and \( m \in L^\infty(\Omega) \). If \( h : \mathbb{R}^d \to \mathbb{R} \) satisfies

\[
- \Delta h = c_d \left( \sum_{p \in \Lambda} \delta_p - m \right) \quad \text{in } \Omega
\]

in the sense of distributions, where \( \Lambda \) is a discrete set of \( \mathbb{R}^d \), we define for any nonnegative function \( \chi \) (typically a cut-off function) the quantity \( W(\nabla h, \chi) \) to be

\[
W(\nabla h, \chi) := \lim_{\eta \to 0} \left( \int_{\Omega \cup \bigcup_{p \in \Lambda} B(p, \eta)} \chi |\nabla h|^2 - c_d g(\eta) \sum_{p \in \Lambda} \chi(p) \right). \tag{3.10}
\]

This is a way of computing \( \int |\nabla h|^2 \) in a “renormalized” fashion, by cutting small holes around the points where \( |\nabla h|^2 \) diverges and subtracting off the expected divergence \( c_d g(\eta) \) corresponding to each point. The name renormalized energy originates in [BBH] where this way of computing was first introduced, in dimension 2, in the context of Ginzburg-Landau vortices. We may check the very existence of the limit in (3.10) by writing that near each point \( p \in \Lambda \), we have

\[
h(x) = c_d g(x - p) + f(x) \quad \text{with } f \text{ of class } C^1.
\]

Thus

\[
\int_{B(p, \eta)} |\nabla h|^2 \approx c_d^2 \int_{B(p, \eta)} |\nabla g|^2 \approx c_d g(\eta) \quad \text{as } \eta \to 0,
\]

which is the contribution that appears in the second term of the right-hand side in (3.10), the precise computation will be done below.

The following lemma connects this quantity with that of the splitting formula (3.4).
Lemma 3.2. Let $d \geq 2$, and $\mu_0$ a measure with a bounded density. For any configuration of distinct points $(x_1, \ldots, x_n) \in (\mathbb{R}^d)^n$, $h_n$ being defined in (3.5), the following identity holds:

$$\int_{\Delta^c} g(x-y) d(\nu_n - n\mu_0)(x) d(\nu_n - n\mu_0)(y) = \frac{1}{c_d} W(\nabla h_n, 1_{\mathbb{R}^d}). \quad (3.11)$$

**Proof.** To prove the lemma, let us compute $W(\nabla h_n, 1_{\mathbb{R}^d})$. We may choose $R$ large enough so that all the points are contained in the ball $B_R = B(0, R)$. In the rest of this text we will use the notation $\nu$ for the outer unit normal to a set. This should not be confused some measures $\nu$ that will appear but we believe the formulas will rarely be ambiguous in that respect. We will thus also denote by $\frac{\partial h}{\partial \nu}$ the normal component of $\nabla h$ on the boundary of a set.

Let $\eta > 0$ be given. By Green’s formula, we have

$$\int_{B_R \setminus \cup_{i=1}^n B(x_i, \eta)} |\nabla h_n|^2 = \int_{\partial B_R} h_n \frac{\partial h_n}{\partial \nu} - \sum_{i=1}^n \int_{\partial B(x_i, \eta)} h_n \frac{\partial h_n}{\partial \nu} - \int_{B_R \setminus \cup_{i=1}^n B(x_i, \eta)} h_n \Delta h_n$$

$$= \int_{\partial B_R} h_n \frac{\partial h_n}{\partial \nu} - \sum_{i=1}^n \int_{\partial B(x_i, \eta)} h_n \frac{\partial h_n}{\partial \nu} - n c_d \int_{B_R \setminus \cup_{i=1}^n B(x_i, \eta)} h_n d\mu_0. \quad (3.12)$$

Here we have used (3.6) and the fact that $\nu_n = 0$ outside of $\cup B(x_i, \eta)$. The normal $\nu$ to each ball $B(x_i, \eta)$ is pointing outwards. Let us now use (temporarily) the notation $h_i^\ast(x) = h_n(x) - g(x - x_i)$ (for the potential generated by the distribution bereft of the point $x_i$). The function $h_i^\ast$ is regular near $x_i$. Using this fact, we may write

$$\int_{\partial B(x_i, \eta)} h_n \frac{\partial h_n}{\partial \nu} = \int_{\partial B(x_i, \eta)} g(-x_i) \frac{\partial h_n}{\partial \nu} + \int_{\partial B(x_i, \eta)} h_i^\ast \frac{\partial h_n}{\partial \nu}$$

$$= g(\eta) \int_{\partial B(x_i, \eta)} \frac{\partial h_n}{\partial \nu} + h_i^\ast(x_i) \int_{\partial B(x_i, \eta)} \left( \frac{\partial h_n}{\partial \nu} + o(1) \right). \quad (3.13)$$

We are left with computing

$$\int_{\partial B(x_i, \eta)} \frac{\partial h_n}{\partial \nu} = \int_{B(x_i, \eta)} \Delta h_n = -c_d \int_{B(x_i, \eta)} (\nu_n - n\mu_0) = -c_d + O(\eta^d) \quad \text{as } \eta \to 0. \quad (3.14)$$

Here we used Green’s formula again, (3.6), the fact that for $\eta$ small enough $B(x_i, \eta)$ contains no other point of the configuration, and the fact that $\mu_0$ has a bounded density. The $O(\eta^d)$ above depends on $n$ and on that bound on the density. It thus follows that

$$- \sum_{i=1}^n \int_{\partial B(x_i, \eta)} h_n \frac{\partial h_n}{\partial \nu} = nc_d g(\eta) + c_d \sum_{i=1}^n h_i^\ast(x_i) + o(1) \quad (3.15)$$
using the obvious fact that by definition of \( g \), \( g(\eta)\eta^d = o_\eta(1) \) as \( \eta \to 0 \). We may insert this into (3.12) while letting \( R \to \infty \). The decay of \( h_n \) and \( \nabla h_n \) mentioned above ensures that
\[
\int_{\partial B_R} h_n \frac{\partial h_n}{\partial \nu} \to 0 \text{ as } R \to +\infty.
\]
We thus obtain
\[
\int_{\mathbb{R}^d \cup B(x_i, \eta)} |\nabla h_n|^2 - n c_d g(\eta) d\mu_0 = c_d \sum_{i=1}^n h^\dagger_n(x_i) - n c_d \int_{\mathbb{R}^d} h_n d\mu_0 + o_\eta(1) \text{ as } \eta \to 0,
\]
where we have used the bound on \( \mu_0 \) and the control of \( h_n \) by \( g \) to show that
\[
\int_{B(x_i, \eta)} h_n d\mu_0 \leq O(\int_{B(0, \eta)} g(x) dx) = O_\eta(1) \text{ by definition of } g.
\]
By definition of \( W \) in (3.10), letting \( \eta \to 0^+ \), this is exactly
\[
W(\nabla h_n, 1_{\mathbb{R}^d}) = c_d \sum_{i=1}^n h^\dagger_n(x_i) - n c_d \int_{\mathbb{R}^d} h_n d\mu_0.
\]
(3.16)

Now, from the definitions it is easily seen that
\[
h^\dagger_n(x_i) = \int_{\mathbb{R}^d \setminus \{x_i\}} g(x - y) d(\nu_n - n\mu_0)(y),
\]
from which it follows that
\[
\iint_{\triangle^n} g(x - y) d(\nu_n - n\mu_0)(x) d(\nu_n - n\mu_0)(y)
\]
\[
= \sum_{i=1}^n \int_{\mathbb{R}^d \setminus \{x_i\}} g(x_i - y) d(\nu_n - n\mu_0)(y) - n \int_{\mathbb{R}^d} h_n d\mu_0
\]
\[
= \sum_{i=1}^n h^\dagger_n(x_i) - n \int h_n d\mu_0.
\]
In view of (3.16), we conclude that (3.11) holds. \( \Box \)

**Remark 3.3.** When examining carefully the proof, in particular near (3.15) and below, we see that one does not really need \( \mu_0 \) to have bounded density; instead it suffices to assume that for any \( x \), \( \mu_0(B(x, \eta)) \) grows no faster than \( \eta^\alpha \) with \( \alpha > d - 2 \) so that \( g(\eta)\eta^\alpha \to 0 \) as \( \eta \to 0 \), and that \( \int_{B(x, \eta)} g(y - x) d\mu_0(y) \to 0 \) as \( \eta \to 0 \), which encompasses a whole class of singular measures, such as measures supported on a rectifiable \((d - 1)\)-dimensional set. Accordingly, Definition 3.1 can be extended with such measures replacing the density \( m \).

Before stating the final formula, we want to blow-up at a scale where the points are well-separated. The convergence of the empirical measure of the \( n \) points to a fixed compactly supported measure suggests that there are typically \( n \) points in a bounded domain, so that the distance between two points should be of order \( n^{1/d} \). To get a \( O(1) \) distance, one thus has to change the scale by a factor \( n^{1/d} \). We will use a ‘ (prime) symbol to denote the blown-up quantities:
\[
x' := n^{1/d} x \text{ for all } x \in \mathbb{R}^d, \text{ in particular } x'_i = n^{1/d} x_i.
\]
Let us also define:

\[ h'_n := g * \left( \sum_{i=1}^{n} \delta_{x'_i} - \mu'_0 \right), \]

so that \( h'_n(x') = n^{2/d-1} h_n(x) \) \( (3.18) \)

where \( \mu'_0 \) is the blown-up measure associated to \( \mu_0 \):

\[ dp'_0(x') = m_0(x) \, dx' = m_0(x' n^{-1/d}) \, dx'. \]

Note that \( \mu'_0 \) has total mass \( n \).

The function \( h'_n \) is the potential generated by the blown-up distribution \( \sum_{i=1}^{n} \delta_{x'_i} - \mu'_0 \) of support \( \Sigma' = n^{1/d} \Sigma \). These modifications affect the expression \( W \) in the following way: since, by a change of variables,

\[ \int_{\mathbb{R}^{d} \setminus \cup_{i=1}^{n} B(x'_i, \eta_n)} |\nabla h'_n|^2 = n^{1-2/d} \int_{\mathbb{R}^{d} \setminus \cup_{i=1}^{n} B(x'_i, n^{1/d} \eta)} |\nabla h_n|^2, \]

we also have that

\[ W(\nabla h_n, 1_{\mathbb{R}^{d}}) = \lim_{\eta_n \to 0} \left[ n^{1-2/d} \int_{\mathbb{R}^{d} \setminus \cup_{i=1}^{n} B(x'_i, n^{1/d} \eta)} |\nabla h'_n|^2 - nc_d g'(\eta') \right] \]

\[ = \lim_{\eta' \to 0} \left[ n^{1-2/d} \int_{\mathbb{R}^{d} \setminus \cup_{i=1}^{n} B(x'_i, \eta')} |\nabla h'_n|^2 - nc_d g'(\eta' n^{-1/d}) \right]. \] \( (3.19) \)

For \( d = 2 \), one has \( g'(\eta' n^{-1/d}) = -\frac{1}{2} \log n + g(\eta') \) whereas for \( d \geq 3 \), \( g(\eta' n^{-1/d}) = n^{2/d} g(\eta') \). Consequently we get the following formulae for the change of scale on the expression of the precursor to the renormalized energy:

\[ W(\nabla h_n, 1_{\mathbb{R}^{d}}) = W(\nabla h'_n, 1_{\mathbb{R}^{d}}) - \frac{cd}{2} n \log n \quad \text{if} \quad d = 2 \]

\[ W(\nabla h_n, 1_{\mathbb{R}^{d}}) = n^{1-2/d} W(\nabla h'_n, 1_{\mathbb{R}^{d}}) \quad \text{if} \quad d \geq 3 \] \( (3.20) \)

Equation \( 3.3 \) and Lemma \( 3.2 \) together with \( 3.20 \) give

**Proposition 3.4** (Splitting formula). *Let \( V \) be such that a unique equilibrium measure \( \mu_0 \) exists, and \( \mu_0 \) has an \( L^\infty \) density (or is as in Remark 3.3). For any \( n \geq 1 \), for any configuration of distinct points \( x_1, \ldots, x_n \) in \( \mathbb{R}^{d}, d \geq 2 \), the following identity holds, with \( \zeta \) as in \( (2.36) \) and \( h'_n \) as in \( (3.18) \):

\[
H_n(x_1, \ldots, x_n) = n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) + \left( -\frac{n}{2} \log n \right) 1_{d=2} + \frac{n^{1-2/d}}{c_d} W(\nabla h'_n, 1_{\mathbb{R}^{d}}). 
\]

We emphasize that here again there is no error term, it is an equality for every \( n \) and every configuration. This formula was first established in [SS7] in dimension \( d = 2 \), and generalized (with the same proof) to higher dimension in [RouSc].
3.3 The case $d = 1$

Since $\zeta$ plays no other role than confining the points to $\Sigma$ (the support of $\mu_0$), this formula shows that it suffices to analyze the behavior of $W(\nabla h'_n, 1_{\mathbb{R}^d})$. We will show later that for good configurations (those that do not have too much energy) $W(\nabla h'_n, 1_{\mathbb{R}^d})$ is proportional to $n$, the number of points. We have thus separated orders in the expansion of $H_n$: after the leading order term $n^2 I(\mu_0)$ and an exceptional term $-\frac{1}{2}n \log n$ in dimension 2, a next term of order $n^{2-2/d}$ appears.

The analysis of $\frac{1}{n} W(\nabla h'_n, 1_{\mathbb{R}^d})$ as $n$ goes to infinity will be the object of Chapters 4 and 6.

3.3 The case $d = 1$

With our choice of $g(x) = -\log |x|$ in dimension 1, $g$ is no longer the Coulomb kernel, so the formal computation (3.8) does not work. However, $g$ is the kernel of the half-Laplacian, and it is known that the half-Laplacian can be made to correspond to the Laplacian by adding one extra space dimension. In other words, we should imbed the one-dimensional space $\mathbb{R}$ into the two-dimensional space $\mathbb{R}^2$ and consider the harmonic extension of $h_n$, defined in (3.5), to the whole plane. That extension will solve an appropriate Laplace equation, and we will reduce dimension 1 to a special case of dimension 2. This is the approach proposed in [SSS].

Let us now get more specific. Let us consider $\mu_0$ the one-dimensional equilibrium measure associated to $I$ (with $g = -\log$) as in Theorem 2.1, and assume it has an $L^\infty$ density $m_0(x)$ with respect to the (one-dimensional) Lebesgue measure (this happens for example when $V(x) = x^2$, the corresponding equilibrium measure being the semi-circle law, cf. Example 2 in Chapter 2). We may now view $\mu_0$ as a singular measure on $\mathbb{R}^2$ by setting

$$d\mu_0(x, y) = m_0(x)\delta_{\mathbb{R}}$$

(3.21)

where $\delta_{\mathbb{R}}$ is the measure of length on the real axis. More precisely, we define $\delta_{\mathbb{R}}$ by its action against smooth test functions $\varphi \in C^0_0(\mathbb{R}^2)$ by

$$\langle \delta_{\mathbb{R}}, \varphi \rangle := \int_{\mathbb{R}} \varphi(x, 0) dx$$

(3.22)

which makes $\delta_{\mathbb{R}}$ a Radon measure on $\mathbb{R}^2$, supported on the real axis. Given $x_1, \ldots, x_n \in \mathbb{R}$, let us also identify them with the points $(x_1, 0), \ldots, (x_n, 0)$ on the real axis of $\mathbb{R}^2$. We may then define the potential $h_n$ on $\mathbb{R}^2$ by

$$h_n = g * \left( \sum_{i=1}^{n} \delta_{(x_i, 0)} - nm_0 \delta_{\mathbb{R}} \right).$$

with $g(x) = -\log |x|$ in $\mathbb{R}^2$, which is nothing else than the harmonic extension to $\mathbb{R}^2$, away from the real axis, of the potential $h_n$ defined in dimension 1 in (3.5).
Viewed as a function in $\mathbb{R}^2$, $h_n$ solves

$$-\Delta h_n = c_2 \left( \sum_{i=1}^{n} \delta_{(x_i,0)} - nm_0 \delta_{\mathbb{R}} \right) \quad \text{in } \mathbb{R}^2$$

which is now a local equation, in contrast with (3.7). We then observe that, letting $\nu_n = \sum_{i=1}^{n} \delta_{(x_i,0)}$, we may also write

$$H_n(x_1, \ldots, x_n) = \int_{\Delta \subset \mathbb{R}^2 \times \mathbb{R}^2} g(x - y) d\nu_n(x) d\nu_n(y) + \int_{\mathbb{R}^2} V d\nu_n,$$

where $g$ is always $-\log |x|$, and $V$ is arbitrarily extended to $\mathbb{R}^2$. This is thus formally the same as in dimension 2, so returning to the setting of Section 3.1 an viewing $\nu_n$ and $\mu_0$ as measures in $\mathbb{R}^2$ thanks to (3.21), we may carry on with the proof of the splitting formula as in the case $d = 2$. We need to use the result of Remark 3.3, which applies because precisely $\mu_0$ is a singular measure but absolutely continuous and with bounded density with respect to the Hausdorff measure on the real axis. The proof then goes through with no other change and yields the same first splitting formula

$$H_n(x_1, \ldots, x_n) = n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) + \frac{1}{c_2} W(\nabla h_n, 1_{\mathbb{R}^2}). \quad (3.23)$$

Continuing on with the blow-up procedure, the natural change of scale is then

$$x' = nx \quad \mu'_0(x') = m_0(x' n^{-1}) \delta_{\mathbb{R}} \quad h'_n(x') = g * \left( \sum_{i=1}^{n} \delta_{(x'_i,0)} - \mu'_0 \right), \quad (3.24)$$

and thus in (3.19) we obtain the term $g(\eta' n^{-1}) = g(\eta') - \log n$ which yields the formula for the change of scales

$$W(\nabla h_n, 1_{\mathbb{R}^2}) = W(\nabla h'_n, 1_{\mathbb{R}^2}) - c_2 n \log n.$$

We conclude with the following splitting formula for $d = 1$ :

**Proposition 3.5.** Let $d = 1$ and $V$ be such that a unique equilibrium measure $\mu_0$ minimizing $I$ exists, and $\mu_0$ has an $L^\infty$ density. Then, for any $n \geq 1$ and any configuration of distinct points $x_1, \ldots, x_n \in \mathbb{R}$, the following identity holds, with $\zeta$ as in (2.36) and $h'_n$ as in (3.24) :

$$H_n(x_1, \ldots, x_n) = n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) - n \log n + \frac{1}{c_2} W(\nabla h'_n, 1_{\mathbb{R}^2}).$$

There remains to understand the term $W(\nabla h'_n, 1_{\mathbb{R}^2})$ as before, except for the particularity that $h'_n$ solves :

$$- \Delta h'_n = c_2 \left( \sum_{i=1}^{n} \delta_{(x'_i,0)} - m_0(x) \delta_{\mathbb{R}} \right) \quad (3.25)$$
with the extra $\delta_R$ term. In the rest of these notes, we will not expand much on
the one-dimensional case. The idea is that with the above transformations, it can essentially be treated almost like the two-dimensional case. The interested reader can refer to [SS8].

### 3.4 Lower bound for the renormalized energy pre-cursor in the case of well-separated points

One of our essential tasks will be to obtain general lower bounds for quantities of the form $W(\nabla h, \chi)$, where $h$ is a potential generated by point charges and a background charge. Such lower bounds are in general delicate, except if one knows a priori that the points are separated by (at least) a fixed distance. We present the lower bound in this case, it is simple yet instructive and will be used later. The elementary proof is a baby version of the “ball construction” method introduced to obtain lower bounds for Ginzburg-Landau functionals by Jerrard [Je], Sandier [Sa], on which we will say more in Chapter 8.

**Lemma 3.6** (Lower bound for well separated points). Let $d \geq 2$ and assume that $h$ solves a relation of the form

$$-\Delta h = c_d \left( \sum_{p \in \Lambda} \delta_p - \mu(x) \right)$$

in an open set $\Omega \subset \mathbb{R}^d \quad (3.26)$

in the sense of distributions, for some discrete set $\Lambda$ and a measure $\mu$ as in Remark 3.3. Assume that the points are well-separated in the sense that for some $r_0 > 0$,

$$\min \left( \min_{p \neq p' \in \Lambda} |p - p'|, \min_{p \in \Lambda} \text{dist}(p, \partial \Omega) \right) \geq 2r_0 > 0. \quad (3.27)$$

Then we have

$$W(\nabla h, 1_\Omega) \geq -C \#(\Lambda \cap \Omega) \quad (3.28)$$

where $C > 0$ depends only on $\mu$, $r_0$ and $d$.

**Proof.** The proof consists in proving lower bounds on disjoint annuli centered at the points of $\Lambda$ via the Cauchy-Schwarz inequality. For any $p \in \Lambda$ and any $\eta < r_0$ we may write that

$$\int_{B(p,r_0) \setminus B(p,\eta)} |\nabla h|^2 = \int_{\eta}^{r_0} \left( \int_{\partial B(p,t)} |\nabla h|^2 \right) dt$$

then use that $|\nabla h|^2 \geq |\frac{\partial h}{\partial \nu}|^2$, and apply the Cauchy-Schwarz inequality on each sphere $\partial B(p,t)$ to get that

$$\int_{B(p,r_0) \setminus B(p,\eta)} |\nabla h|^2 \geq \int_{\eta}^{r_0} \left( \int_{\partial B(p,t)} \frac{\partial h}{\partial \nu} \right)^2 \frac{1}{|S^{d-1}| \nu^{d-1}} \, dt. \quad (3.29)$$
By Green’s formula and integration of (3.26), we have
\[ \int_{\partial B(p,t)} \frac{\partial h}{\partial \nu} = \int_{B(p,t)} \Delta h = -c_d + O(\mu(B(p,t))), \]
because for \( t \leq r_0 \), \( B(p,t) \cap \Lambda = \{ p \} \) by (3.27). Since we assume that \( \mu(B(p,t)) \leq C_0 t^\alpha \) with \( \alpha > d - 2 \), inserting this into (3.29) yields
\[ \int_{B(p,r_0) \setminus B(p,\eta)} |\nabla h|^2 \geq \int_\eta^{r_0} \frac{c_d^2 - Ct^\alpha}{|\mathbb{S}^{d-1}|d-1} dt \tag{3.30} \]
for some \( C \) depending only on \( C_0 \) (hence on \( \mu \)) and on \( d \). Next, we may compute explicitly, for \( d \geq 3 \),
\[ \int_\eta^{r_0} \frac{c_d^2}{|\mathbb{S}^{d-1}|d-1} dt = \frac{c_d^2}{(2 - d)|\mathbb{S}^{d-1}|} \left( c_0^{2-d} - \eta^{2-d} \right) = c_d(g(\eta) - g(r_0)) \tag{3.31} \]
where we have used that \( c_d = (d - 2)|\mathbb{S}^{d-1}| \). In dimension \( d = 2 \) where \( c_2 = 2\pi \), we can check that the same conclusion holds.

On the other hand, the second term in (3.30) is easily bounded:
\[ \int_\eta^{r_0} \frac{Ct^\alpha}{|\mathbb{S}^{d-1}|d-1} dt \leq C(\mu, r_0, d), \tag{3.32} \]
since \( \alpha > d - 2 \), with \( C \) a constant depending only on \( C_0, r_0 \) and \( d \). Combining (3.30), (3.31) and (3.32) we eventually get that
\[ \int_{B(p,r_0) \setminus B(p,\eta)} |\nabla h|^2 \geq c_d(g(\eta) - g(r_0)) - C(\mu, r_0, d). \]
Since the balls \( B(p,r_0) \) are disjoint by assumption, we may add all their contributions and compare with the definition (3.10) to get (3.28). \( \square \)

**Remark 3.7.** Assuming (3.26) holds in a domain \( \Omega \), if \( B(p,\eta) \) contains \( N_p \) points, the same calculation as in (3.29)–(3.31) gives
\[ \int_{B(p,r_0) \setminus B(p,\eta)} |\nabla h|^2 \geq c_d N_p^2 g(\eta) - C. \]
This way we see that each point \( p \) with multiplicity \( N_p \) costs \( c_d N_p^2 g(\eta) \) (to leading order) in the energy \( \int |\nabla h|^2 \), and this is quadratic in \( N_p \). Thus, when subtracting off \( c_d \chi(p)g(\eta) \) in the definition of \( W(\nabla h, \chi) \), if there is one multiple point i.e. \( N_p \geq 2 \) we are not subtracting enough and the limit defining \( W(\nabla h, \chi) \) gives \( +\infty \).
4 Lower bounds by smearing out the charges

In this chapter we introduce a first way of obtaining lower bounds for quantities involving $W$, based on “smearing out” (in a sense that will be made precise) the singular charges. This method relies on this smearing out idea that has been used for a long time, in particular to prove Onsager’s lemma [Ons], a tool which has been much used in the proof of stability of matter in quantum mechanics (see [LieOx, LieSei1] and references therein). In this chapter we follow closely [RouSe]. We will restrict to $d \geq 2$, with the understanding that the case $d = 1$ can be treated as indicated in the previous chapter. The complete definitions of the renormalized energy itself will be given in the next chapter.

4.1 Definition of the charge smearing out and Newton’s theorem

Let us fix a nonnegative function $\rho$, radial (this is an important assumption), supported in the unit ball $B(0, 1)$ of $\mathbb{R}^d$ and of integral 1. A simple example is to take $\rho = \frac{1_{B(0,1)}}{|B(0,1)|}$. We then define, for any $x \in \mathbb{R}^d$ and for any $\eta > 0$, the function $\delta^{(\eta)}_x$:

$$\delta^{(\eta)}_x = \frac{1}{\eta^d} \rho \left( x + \frac{\cdot}{\eta} \right) = \frac{1}{\eta^d} \rho \left( \frac{\cdot}{\eta} \right) * \delta_x. \quad (4.1)$$

The function $\delta^{(\eta)}_x$ is supported in $B(x, \eta)$, and its integral is 1. It should be seen as an approximation of the Dirac mass at $x$, obtained by smearing it out radially at the scale $\eta$, in other words a smeared out Dirac mass.

The celebrated Newton’s theorem (see e.g. [LiLo, Theorem 9.7]) allows one to compare the effect of a charge and the one of its smeared out alter ego with respect to Coulomb interactions.

**Theorem 4.1** (Newton’s theorem). The gravitational potential created by a singular charge at the origin and that created by a radial distribution of same mass contained in a ball centered at the origin, coincide outside the ball. That is, with our notation,

$$g * \delta^{(\eta)}_0 = g \text{ outside } B(0, \eta). \quad (4.2)$$

Moreover, it holds that

$$g * \delta^{(\eta)}_0 \leq g \text{ everywhere in } \mathbb{R}^d. \quad (4.3)$$
Proof. For every \( x \in B(0, \eta)^c \), we know that \( g \) is harmonic in \( B(x, t) \) for every \( t < \eta \). Thus, by the mean-value property of harmonic functions, we have

\[
\int_{\partial B(x, t)} g = g(x),
\]

and it follows, using the radial character of \( \rho \), that

\[
g \ast \delta_0^{(\eta)}(x) = \frac{1}{\eta^d} \int_{B(0, \eta)} g(x-y) \rho \left( \frac{y}{\eta} \right) dy = \frac{1}{\eta^d} \int_0^\eta \rho \left( \frac{t}{\eta} \right) \left[ \int_{\partial B(0, t)} g(x-y) dy \right] dt
\]

\[
= \frac{1}{\eta^d} \int_0^\eta \rho \left( \frac{t}{\eta} \right) |\partial B(0, t)| g(x) dt
\]

\[
= g(x) \int_0^1 \rho(t) |\partial B(0, t)| dt = g(x) \int_{B(0,1)} \rho = g(x). \quad (4.4)
\]

Hence \( g \ast \delta_0^{(\eta)}(x) = g(x) \) for any \( x \) outside \( B(0, \eta) \), as claimed. Moreover, if \( x \) is an arbitrary point, the second line in (4.4) becomes an inequality for any \( x \) by the fact that \( g \) is always super-harmonic, hence

\[
g \ast \delta_0^{(\eta)}(x) \leq \frac{1}{\eta^d} \int_0^\eta \rho \left( \frac{t}{\eta} \right) |\partial B(0, t)| g(x) dt = g(x).
\]

Corollary 4.1. For any \( \eta > 0 \), the radial function \( f_\eta = g \ast \delta_0^{(\eta)} - g \) is nonpositive and solution to

\[
\begin{cases}
-\Delta f_\eta = c_d \left( \delta_0^{(\eta)} - \delta_0 \right) & \text{in } \mathbb{R}^d \\
f_\eta \equiv 0 & \text{in } \mathbb{R}^d \setminus B(0, \eta).
\end{cases}
\quad (4.5)
\]

The function \( f_\eta \) will be used in this chapter and the next. We will also sometimes abuse notation by considering \( f_\eta \) as a function of the real variable corresponding to the distance to the origin.

### 4.2 Coulomb interactions and Onsager’s lemma

Let us start by introducing the notation for the Coulomb interaction that we will use in the sequel, and which is most commonly used in the literature. For two charge distributions (bounded Radon measures) \( \mu_1 \) and \( \mu_2 \), one sets

\[
D(\mu_1, \mu_2) := \iint_{\mathbb{R}^d \times \mathbb{R}^d} g(x-y) \, d\mu_1(x) \, d\mu_2(y). \quad (4.6)
\]
This is formally linked to the (electrostatic) potentials $h_{\mu_1} = g*\mu_1$, and $h_{\mu_2} = g*\mu_2$
that the distributions generate via the formula (obtained by formal integration by
parts and using that $-\Delta h_{\mu_i} = c_d \mu_i$)
\[
D(\mu_1, \mu_2) = \int_{\mathbb{R}^d} \mu_1 h_{\mu_2} = \int_{\mathbb{R}^d} h_{\mu_1} = \frac{1}{c_d} \int_{\mathbb{R}^d} \nabla h_{\mu_1} \cdot \nabla h_{\mu_2}
\]
which is however not always true, particularly in dimension 2. It is easily seen
that $D(f, f) \geq 0$ and vanishes only for $f = 0$. For various other properties of $D,$
we refer to [SaTo] for $d = 2$ and [LiLo, Sec. 9.8.] for $d \geq 3.$

The following lemma easily follows from Newton’s theorem:

**Lemma 4.2 (Onsager’s lemma).**
For any $n$, any $x_1, \ldots, x_n \in \mathbb{R}^d$, any nonnegative distribution $\mu$ such that $\int_{\mathbb{R}^d} \mu = n$, and any $\ell > 0$, we have
\[
\sum_{i \neq j} g(x_i - x_j) \geq D\left(\mu - \sum_{i=1}^n \delta_{x_i}^{(\ell)}, \mu - \sum_{i=1}^n \delta_{x_i}^{(\ell)} \right) - D(\mu, \mu) + 2 \sum_{i=1}^n D(\mu, \delta_{x_i}^{(\ell)}) - nD(\delta_0^{(\ell)}, \delta_0^{(\ell)})
\]
with equality if $\min_{i \neq j} |x_i - x_j| \geq 2\ell.$ In this equation, the left-hand side is the
true total Coulomb interaction of the points and the right-hand side is its “smeared
out” version.

This lemma allows to compare the Coulomb interaction between $n$ points with
the interaction of these points and a reference measure $\mu$, and such a comparison
is perfectly suited to our study.

**Proof.** From Newton’s theorem, we have
\[
\int \left( g * \delta_{x_i}^{(\ell)} \right) \delta_{x_j}^{(\ell)} \leq \int \left( g * \delta_{x_j}^{(\ell)} \right) \delta_{x_i}^{(\ell)} \leq \int \left( g * \delta_{x_j}^{(\ell)} \right) \delta_{x_i}^{(\ell)}
\]
It thus follows that
\[
\sum_{i \neq j} g(x_i - x_j) \geq \sum_{i \neq j} D(\delta_{x_i}^{(\ell)}, \delta_{x_j}^{(\ell)})
\]
with equality if $\min_{i \neq j} |x_i - x_j| \geq 2\ell.$ In this equation, the left-hand side is the
true total Coulomb interaction of the points and the right-hand side is its “smeared
out” version.

We note that this inequality is independent of the mutual distance between
the points, and the points can be arbitrarily close from each other (in particular, the
smeared out charges can overlap). On the other hand, it does not behave nicely if
one considers charges of different signs.
After this application of Newton’s theorem, the lemma reduces to an algebraic identity: expanding by bilinearity, we then write

\[
D \left( \mu - \sum_{i=1}^{n} \delta_{\ell_i}, \mu - \sum_{i=1}^{n} \delta_{\ell_i} \right) = D(\mu, \mu) - 2 \sum_{i=1}^{n} D(\mu, \delta_{\ell_i}) + \sum_{i \neq j} D(\delta_{\ell_i}, \delta_{\ell_j})
\]

and from this relation and (4.9), by translation-invariance of the Coulomb interaction, the lemma easily follows.

Smearing out charges comes with a cost, that we quantify in the following lemma.

**Lemma 4.3** (The cost of smearing out charges).

For any \( \mu \in L^\infty(\mathbb{R}^d) \) and any point \( x \), we have

\[
\left| D \left( \mu, \delta_x - \delta_x \right) \right| \leq C \ell^2 \| \mu \|_{L^\infty},
\]

where \( C > 0 \) depends only on the choice of \( \rho \) and the dimension.

**Proof.** Without loss of generality, we may assume \( x = 0 \), and write

\[
D(\mu, \delta_0 - \delta_0) = \int (g - g * \delta_0)(x) \mu(x) \, dx.
\]

By Newton’s theorem, the function \( g - g * \delta_0 \) is nonnegative and supported in \( B(0, \ell) \). In dimension 3, we may just write that in particular it is smaller than \( g 1_{B(0, \ell)} \), and so we may write

\[
\left| D(\mu, \delta_0 - \delta_0) \right| \leq \| \mu \|_{L^\infty} \int_{B(0, \ell)} \frac{dx}{|x|^{d-2}} \leq C \ell^2 \| \mu \|_{L^\infty}.
\]

In dimension 2, we write instead

\[
\int_{\mathbb{R}^2} |g - g * \delta_0| = \int_{B(0, \ell)} \left| - \log |x| + \frac{1}{\ell^2} \int \log |x - y| \rho(y) \, dy \right| \, dx
\]

\[
= \ell^2 \int_{B(0, 1)} \left| - \log |x'| + \int \log |x' - y'| \rho(y') \, dy' \right| \, dx = C \ell^2
\]

where we have used the changes of variables \( x = \ell x', y = \ell y' \), and we conclude in the same way. \( \square \)
4.3 Splitting lower bound

In this section, we see how to use the smearing out and Onsager’s lemma to deduce lower bounds for the Coulomb gas Hamiltonian, more specifically for the part \( W(\nabla h_n, 1_{\mathbb{R}^d}) \) in the splitting formula of Proposition 3.4.

The idea is simply to apply Lemma 4.2 with the reference measure \( \mu = n\mu_0 \) and the smearing out scale \( \ell = \eta^{1/d} \), and then to apply the same recombinations of terms as in Chapter 3.

We obtain the following intermediate result :

**Lemma 4.4 (First splitting lower bound).** Assume the equilibrium measure \( \mu_0 \) exists and has an \( L^\infty \) density. For any \( n \) and any \( x_1, \ldots, x_n \in \mathbb{R}^d \), we have

\[
H_n(x_1, \ldots, x_n) \geq n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) + D\left( \sum_{i=1}^{n} \delta^{(\ell)}_{x_i} - n\mu_0, \sum_{i=1}^{n} \delta^{(\ell)}_{x_i} - n\mu_0 \right) - nD(\delta^{(\ell)}_{0}, \delta^{(\ell)}_{0}) - Cn^2 \ell^2,
\]

where \( C \) depends only on the dimension and \( \|\mu_0\|_{L^\infty} \). Moreover, equality holds if \( \min_{i \neq j} |x_i - x_j| \geq 2\ell \).

Note that combined with (3.4) and Proposition 3.2 this yields the inequality

\[
W(\nabla h_n, 1_{\mathbb{R}^d}) \geq c_d D\left( \sum_{i=1}^{n} \delta^{(\ell)}_{x_i} - n\mu_0, \sum_{i=1}^{n} \delta^{(\ell)}_{x_i} - n\mu_0 \right) - nc_d D(\delta^{(\ell)}_{0}, \delta^{(\ell)}_{0}) - Cn^2 \ell^2.
\]

In the right-hand side appear an error term \(-Cn^2 \ell^2\) and a quantity which we will be able to bound from below : it is the total Coulomb interaction between the smeared out charges \( \delta^{(\ell)}_{x_i} \) and the “background charge” \(-n\mu_0\), from which the self-interaction of the charges \( c_d D(\delta^{(\ell)}_{0}, \delta^{(\ell)}_{0}) \) is subtracted. In spirit, unsurprisingly, it is equivalent to \( W(\nabla h_n, 1_{\mathbb{R}^d}) \), but the method for “renormalizing”, i.e. removing the infinite self-interactions of the singular charges is different : it uses the smearing out method instead of cutting out small holes around the Dirac charges. In addition, we will see later that points in minimizing configurations are well-separated so that there is equality in (4.11) and (4.12), and no information is lost in the end. In such a situation, the two methods for “renormalizing” will be shown to be equivalent.

**Proof.** As announced, we apply Lemma 4.2 with the reference measure \( \mu = n\mu_0 \) and the smearing out scale \( \ell = \eta^{1/d} \), and we obtain

\[
H_n(x_1, \ldots, x_n) \geq D\left( n\mu_0 - \sum_{i=1}^{n} \delta^{(\ell)}_{x_i}, n\mu_0 - \sum_{i=1}^{n} \delta^{(\ell)}_{x_i} \right) - n^2 D(\mu_0, \mu_0) + 2n \sum_{i=1}^{n} D(\mu_0, \delta^{(\ell)}_{x_i}) - nD(\delta^{(\ell)}_{0}, \delta^{(\ell)}_{0}) + n \sum_{i=1}^{n} V(x_i).
\]
In view of Lemma 4.3, we may write
\[ 2n \sum_{i=1}^{n} D(\mu_0, \delta^{(i)}_{x_i}) = 2n \sum_{i=1}^{n} D(\mu_0, \delta_{x_i}) + O(n^2 \|\mu_0\|_{L^\infty}). \]

Meanwhile, we observe that \( D(\mu_0, \delta_{x_i}) = \int g(x_i - y) d\mu_0(y) = h^\mu_0(x_i) \), where \( h^\mu_0 \) was defined in (2.16). Hence we may recombine the terms
\[ 2n \sum_{i=1}^{n} D(\mu_0, \delta_{x_i}) + n \sum_{i=1}^{n} V(x_i) = n \sum_{i=1}^{n} (2h^\mu_0 + V)(x_i) = 2n \sum_{i=1}^{n} (\zeta + c)(x_i), \]
where \( \zeta \) is as in (2.36) and \( c \) as in (2.17). Inserting these relations into (4.13), we are led to
\[ H_n(x_1, \ldots, x_n) \geq D(n\mu_0 - \sum_{i=1}^{n} \delta^{(i)}_{x_i}, n\mu_0 - \sum_{i=1}^{n} \delta^{(i)}_{x_i}) - nD(\delta^{(1)}_{0}, \delta^{(1)}_{0}) + 2n \sum_{i=1}^{n} \zeta(x_i) + 2cn^2 - n^2 D(\mu_0, \mu_0) - Cn^2 \|\mu_0\|_{L^\infty}. \]

Inserting the value for \( c = I(\mu_0) - \frac{1}{2} \int V d\mu_0 \) we find that \( 2cn^2 - n^2 D(\mu_0, \mu_0) = n^2 I(\mu_0) \), and we get the result.

The second and final step is to insert into this result the exact value of the self-interaction of a smeared out charge, and to make the same blow up as in Chapter 3. For the latter, we define
\[ \kappa_d := c_d D(\delta^{(1)}_{0}, \delta^{(1)}_{0}) \quad \text{for } d \geq 3, \quad \kappa_2 := c_2, \quad \gamma_2 := c_2 D(\delta^{(1)}_{0}, \delta^{(1)}_{0}) \quad \text{for } d = 2. \]

The numbers \( \kappa_d, \gamma_2 \), correspond to the self-interaction energy of the smeared out Dirac at scale 1. They depend (only) on the choice of the function \( \rho \) and on the dimension. We will see that our results do not however depend on the choice of \( \rho \). The definition is somewhat different and not symmetric when \( d = 2 \), this is due to the fact that the logarithm behaves differently from power functions under rescaling and is made to ease the formulas below.

The scaled versions of the self-interaction energy easily follow from (4.15) and changing variables:
\[
\begin{cases}
    D(\delta^{(\eta)}_{0}, \delta^{(\eta)}_{0}) = \frac{2\eta}{c_d^2} g(\eta) & \text{if } d \geq 3 \\
    D(\delta^{(\eta)}_{0}, \delta^{(\eta)}_{0}) = g(\eta) + \frac{2\eta}{c_2^2} g(\eta) + \frac{2\eta}{c_2} & \text{if } d = 2
\end{cases}
\]
and we have
\[ D(\delta^{(\ell)}_{0}, \delta^{(\ell)}_{0}) = n^{1-2/d} D(\delta^{(\eta)}_{0}, \delta^{(\eta)}_{0}). \]

For the blow-up, let us recall we had defined the potentials
\[ h_n = g * \left( \sum_{i=1}^{n} \delta_{x_i} - n\mu_0 \right) \quad \text{and} \quad h'_n(x') = g * \left( \sum_{i=1}^{n} \delta_{x'_i} - \mu_0 \right) = n^{2/d-1} h_n(x). \]
We now define the potentials generated by the smeared out charges:

\[ h_{n,\ell} = g \left( \sum_{i=1}^{n} \delta_{x_i} - n \mu_0 \right) \quad \text{and} \quad h'_{n,\eta} = g \left( \sum_{i=1}^{n} \delta_{x'_i} - \mu'_0 \right) \]  

with \( \eta = n^{1/d} \ell \). (4.17)

It will be important for later to observe that \( h_n \) and \( h_{n,\ell} \), and \( h'_n \) and \( h'_{n,\eta} \) respectively, are in one-to-one correspondence via

\[ h_{n,\ell} = h_n + \sum_{i=1}^{n} f_{\ell} (x_i) \quad \text{and} \quad h'_{n,\eta} = h'_n + \sum_{i=1}^{n} f_\eta (x'_i), \]  

where the function \( f_\eta \) is as in Corollary 4.1. (4.18)

Using the fact that the total charge is zero and compactly supported, the function \( h_{n,\ell} \) and its gradient decay like \( r^{1-d} \) and \( r^{-d} \) respectively, and so the integration by parts that yields (4.7) is justified, and we may write

\[ D \left( n \mu_0 - \sum_{i=1}^{n} \delta_{x_i}, n \mu_0 - \sum_{i=1}^{n} \delta_{x_i} \right) = \frac{1}{c_d} \int_{\mathbb{R}^d} |\nabla h_{n,\ell}|^2 = \frac{n^{1-2/d}}{c_d} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 \]

Combining these facts with the previous lemma, and using that \( \ell = n^{-1/d} \eta \), we get to the following conclusion:

**Proposition 4.5** (Splitting lower bound [RouSe]). Assume the equilibrium measure \( \mu_0 \) exists and has a bounded density. For every \( n \geq 1 \), and for all configurations of points \( x_1, \ldots, x_n \) in \( \mathbb{R}^d \), the following inequality holds:

\[ H_n(x_1, \ldots, x_n) \geq n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) - \left( \frac{n}{2} \log n \right) 1_{d=2} \]

\[ + \frac{n^{2-2/d}}{c_d} \left( \frac{1}{n} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) - C \eta^2 \right), \]  

where \( h'_{n,\eta} \) is as in (4.17) and \( C > 0 \) depends only on the dimension and on \( \|\mu_0\|_{L^\infty} \). Moreover, equality holds if \( \min_{i \neq j} |x'_i - x'_j| \geq 2\eta \).

Comparing again with the exact splitting formula of Proposition 3.4, this is equivalent to writing that

\[ W(\nabla h'_{n,1}, 1_{\mathbb{R}^d}) \geq n^{2-2/d} \left( \frac{1}{n} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) - C \eta^2 \right). \]  

In this chapter, we have worked and obtained some results with fixed \( n \) and \( \eta \). In the sequel, we will focus on letting \( \eta \to 0 \) and \( n \to \infty \).
4.4 Consequences

The result of Proposition 4.5, for fixed \( n \) and \( \eta \), already yields some easy consequences, such as a lower bound for \( H_n \). Indeed, taking \( \eta = 1 \), and using the fact that \( \int |\nabla h'_{n,\eta}|^2 \geq 0 \), we get as a corollary

**Corollary 4.6** (An easy lower bound for \( H_n \)). Under the same assumptions, we have

\[
H_n(x_1, \ldots, x_n) \geq n^2 I(\mu_0) + 2n \sum_{i=1}^{n} \zeta(x_i) - \left( \frac{n}{2} \log n \right) 1_{d=2} - Cn^{2-2/d} \tag{4.21}
\]

where the constant \( C \) depends only on \( \|\mu_0\|_{L^\infty} \) and the dimension.

Another way of stating this is that \( W(\nabla h'_n, 1_{\mathbb{R}^d}) \geq -C n^{2-2/d} \) and we see that the next order term in the expansion of \( H_n \) can indeed be expected to be of order \( n^{2-2/d} \) — at least it is bounded below by it.

For illustration, let us show how this lower bound easily translates into an upper bound for the partition function \( Z_{n,\beta} \) (defined in (2.59)) in the case with temperature.

**Corollary 4.7** (An easy upper bound for the partition function). Assume that \( V \) is continuous and satisfies (A3)–(A4) (See (2.62) in Section 2.3.) Assume that \( \mu_0 \) has an \( L^\infty \) density. Then for all \( \beta > 0 \), and for \( n \) large enough, we have

\[
\log Z_{n,\beta} \leq \beta n^2 I(\mu_0) + \left( \frac{\beta}{4} n \log n \right) 1_{d=2} + C \beta n^{2-2/d} + Cn
\]

where \( C \) depends only on \( \mu_0 \) and the dimension.

To prove this, let us state a lemma that we will use repeatedly and that exploits assumption (A4).

**Lemma 4.8.** Assume that \( V \) is continuous and satisfies (A3)–(A4). For any \( \lambda > 0 \) we have

\[
\lim_{n \to +\infty} \left( \int_{(\mathbb{R}^d)^n} e^{-\lambda \beta n \sum_{i=1}^{n} \zeta(x_i)} dx_1 \ldots dx_n \right)^{1/n} = |\omega| \tag{4.22}
\]

where \( \omega = \{ \zeta = 0 \} \), uniformly in \( \beta \in [\beta_0, +\infty) \), \( \beta_0 > 0 \).

**Proof.** First, by separation of variables, we have

\[
\left( \int_{(\mathbb{R}^d)^n} e^{-\lambda \beta n \sum_{i=1}^{n} \zeta(x_i)} dx_1 \ldots dx_n \right)^{1/n} = \int_{\mathbb{R}^d} e^{-\lambda \beta n \zeta(x)} dx.
\]
Second, we recall that since \(\mu_0\) is a compactly supported probability measure, \(h^{\mu_0}\) must asymptotically behave like \(g(x)\) as \(|x| \to \infty\), thus \(\zeta = h^{\mu_0} + \frac{V}{2} - c\) grows like \(g(x) + \frac{V}{2}\). The assumption (A4) thus ensures that there exists some \(\alpha > 0\) such that \(\int e^{-\alpha \zeta(x)} \, dx < +\infty\) and thus for \(\lambda > 0\), \(\beta > 0\), and \(n\) large enough, \(\int e^{-\lambda \beta n \zeta(x)} \, dx < +\infty\). Moreover, by definition of \(\omega\) (cf. (2.37)),

\[
e^{-\lambda \beta n \zeta} \to 1_\omega \quad \text{as } n \to +\infty
\]

pointwise, and \(\omega\) has finite measure in view of the growth of \(h^{\mu_0}\) and thus of \(\zeta\). Moreover these functions are dominated by \(e^{-\alpha \zeta}\) for \(n\) large enough, which is integrable, so the dominated convergence theorem applies and allows to conclude.

\(\Box\)

**Proof of the corollary.** By definition (2.59) we have

\[
\log Z_{n,\beta} = \log \int e^{-\frac{1}{2} H_n(x_1, \ldots, x_n)} \, dx_1 \ldots dx_n
\]

and inserting (4.21), we are led to

\[
\log Z_{n,\beta} \leq -\frac{\beta}{2} n^2 I(\mu_0) + \left(\frac{\beta}{4} n \log n\right) 1_{d=2} + C \beta n^{2-2/d} + \log \left(\int e^{-n \beta \sum_{i=1}^n \zeta(x)} \, dx_1 \ldots dx_n\right). \quad (4.23)
\]

Using Lemma 4.8 to handle the last term, we deduce that

\[
\log Z_{n,\beta} \leq -\frac{\beta}{2} n^2 I(\mu_0) + \left(\frac{\beta}{4} n \log n\right) 1_{d=2} + C \beta n^{2-2/d} + n(\log |\omega| + o_n(1))
\]

which gives the conclusion.

\(\Box\)

### 4.5 Control of the potential and charge fluctuations via smearing out

In this section, we show how the method of smearing out the charges, even at a fixed scale \(\eta\) (which does not need to go to zero), allows to obtain in a simple way some control on \(\nabla h_n\) itself, and also on the points.

Indeed \(W(\nabla h_n, 1_{\mathbb{R}^d})\) can be controlled via Proposition 3.4 if a suitable upper bound for \(H_n\) is known, but it does not immediately yield a control on \(\nabla h_n\) in \(L^p\) spaces, due to the “renormalization” procedure (cutting out of small holes) in the definition of \(W\); while in contrast a control of \(H_n\), hence on \(W(\nabla h_n, 1_{\mathbb{R}^d})\) gives a control on \(\int |\nabla h_n|^2\) for some fixed \(\eta\) (such as \(\eta = 1\)) via Proposition 4.5 or (4.20), and we can show that this control in turn gives the optimal desired control on \(\nabla h_n\), thanks to the correspondence (4.18). This is the object of the following :
Lemma 4.9 (Control of the potential via smearing out). Assume $\mu_0$ is a measure with an $L^\infty$ density. For any $n$, any $x_1,\ldots,x_n \in \mathbb{R}^d$, any $1 \leq q < \frac{d}{d-1}$, any $0 < \eta < 1$ and any $R > 0$ and $K_R = [-R,R]^d$, denoting $\nu'_n = \sum_{i=1}^n \delta_{x'_i}$, and letting $h'_{n,\eta}$, $h'_n$ be as in (3.5), (4.17), we have

$$\|\nabla h'_n\|_{L^q(K_R)} \leq |K_R|^{1/q-1/2} \|\nabla h'_{n,\eta}\|_{L^2(K_R)} + C_{q,\eta} \nu'_n(K_R+\eta)$$

where $C_{q,\eta}$ depends only on $q$, $\eta$ and $d$ and satisfies $C_{q,\eta} \to 0$ when $\eta \to 0$ at fixed $q$.

Remark 4.10. The reason for the condition $q < \frac{d}{d-1}$ is that near each $x'_i$, $h'_n$ has a singularity in $g(x-x'_i)$, hence $\nabla h'_n$ blows up like $|x-x'_i|^{1-d}$ and this is in $L^q$ if and only if $q < \frac{d}{d-1}$. In other words, while $\nabla h'_{n,\eta}$ belongs to $L^2$, $\nabla h'_n$ belongs at most to such $L^q$ spaces.

Proof. By (4.18), we have

$$\nabla h'_n = \nabla h'_{n,\eta} - \sum_{i=1}^n \nabla f_\eta(x-x_i)$$

and thus

$$\|\nabla h'_n\|_{L^q(K_R)} \leq \|\nabla h'_{n,\eta}\|_{L^q(K_R)} + \nu'_n(K_R+\eta) \|\nabla f_\eta\|_{L^q(\mathbb{R}^d)}$$

where we used that if $x \in K_R$ and $\eta < 1$, then $f_\eta(x-x_i) = 0$ if $x_i \notin (K_R+\eta)^c$. A simple application of Hölder’s inequality then yields

$$\|\nabla h'_{n,\eta}\|_{L^q(K_R)} \leq |K_R|^{1/q-1/2} \|\nabla h'_{n,\eta}\|_{L^2(K_R)}$$

and concludes the proof of the inequality, with $C_{q,\eta} := \|\nabla f_\eta\|_{L^q(\mathbb{R}^d)}$.

In addition, we can note that since

$$-\Delta h'_{n,\eta} = c_d \left( \sum_{i=1}^n \delta^{(\eta)}_{x'_i} - \mu'_0 \right)$$

again for fixed $\eta$, controlling $\nabla h'_{n,\eta}$ in $L^2$ (as we do in Proposition 4.5) gives a control on $\sum_i \delta^{(\eta)}_{x'_i}$ (more precisely it gives a control on $\|\sum_i \delta^{(\eta)}_{x'_i} - \mu'_0\|$ in the Sobolev space $H^{-1}_{\text{loc}}$), which suffices, say, to control the number of points in a given region, since controlling the Dirac masses or the smeared out Dirac masses is not much different. These controls in weak spaces of the fluctuations $\sum_i \delta_{x'_i} - \mu'_0$ can in fact be improved, again via the smeared out charges at fixed scale $\eta$, and we have the following result:
Lemma 4.11 (Controlling charge fluctuations). Assume $\mu_0$ is a measure with an $L^\infty$ density. For any $n$, for any $x_1,\ldots,x_n \in \mathbb{R}^d$, let $\nu'_n = \sum_{i=1}^n \delta_{x'_i}$ and

$$\nabla h'_{n,\eta}(x', R) := \nu'_n(B(x', R)) - \int_{B(x', R)} d\mu'_0.$$ 

Then, $h'_{n,\eta}$ being given by (4.17), for any $0 < \eta < 1$, $R > 2$ and $x' \in \mathbb{R}^d$, we have

$$\int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 \geq C \frac{\Delta(x', R)^2}{R^{d-2}} \min \left(1, \frac{\Delta(x', R)}{R^d}\right),$$  

(4.25)

where $C$ is a constant depending only on $d$.

For the proof, we refer to [RouSe, Lemma 4.6].
5 Definition(s) and properties of the renormalized energy

In the previous chapters, we have seen the definition of the “precursor to the renormalized energy” $W$ (Definition 3.1), and a splitting of the Hamiltonian where lower order terms appear that involve either $W$ (in Proposition 3.4) or the norm of the smeared out potential, in Proposition 4.5. We have worked so far at fixed $n$. The ultimate goal is to study the asymptotic limit of these lower order terms as $n \to +\infty$. When taking the limit, a limiting object will appear, which we call again a renormalized energy. This energy is the total Coulomb interaction energy of an infinite configuration of points in the whole space in a constant neutralizing background. Such a system is called in physics a jellium. This chapter is devoted to the definition(s) of this limiting object itself and the study of some of its properties, before we proceed in the next chapter with deriving it as the $n \to +\infty$ limit. Thus this chapter can be read independently from the rest.

5.1 Motivation and definitions

The goal of this chapter is to define a total Coulomb interaction for an infinite system of discrete point “charges” in a constant neutralizing background of fixed density $m > 0$, related to a potential $h$ that solves (in the sense of distributions)

$$-\Delta h = c_d \left( \sum_{p \in \Lambda} N_p \delta_p - m \right) \text{ in } \mathbb{R}^d, \text{ for } d \geq 2 \quad (5.1)$$

where $\Lambda$ is a discrete set of points in $\mathbb{R}^d$, and $N_p$ are positive integers (the multiplicities of the points), respectively

$$-\Delta h = c_d \left( \sum_{p \in \Lambda} N_p \delta_p - m \delta_R \right) \text{ in } \mathbb{R}^2, \text{ for } d = 1 \quad (5.2)$$

with $\delta_R$ defined in (3.22).

Again, such a system is often called a (classical) jellium in physics. The jellium model was first introduced by [Wi1] in the quantum case, and can be viewed as a toy model for matter: the points charges are then atoms, which interact (via electrostatic forces) with a cloud of electrons of density $m$.

The reason why we need to consider such systems is that in the two previous chapters we dealt with functions $h'_n$ that solved equations of the following type:

$$-\Delta h'_n = c_d \left( \sum_{i=1}^{n} \delta_{x'_i} - \rho'_0 \right) \text{ for } d \geq 2 \quad (5.3)$$

where we had chosen to center the blow-up at, say, the origin 0 (respectively \((3.25)\) for \(d = 1\)). We note that, in that case, the density \(\mu'_0(x')\) equals by definition \(\mu_0(x'n^{-1/d})\), so that, at least if \(\mu_0\) is sufficiently regular, \(\mu'_0(x') \to \mu_0(0)\) pointwise as \(n \to +\infty\), i.e. \(\mu'_0\) converges to a constant. It is constant because \(\mu_0\) varies much slower than the scale of the configuration of discrete points. If we had chosen to blow up around a different point, say \(x_0\), then we would obtain instead the constant \(\mu_0(x_0)\) as the limit. This constant is the local density of the neutralizing background charge. As \(n \to +\infty\), the number of points becomes infinite and they fill up the whole space, so that if we blow-up around an origin which lies in the support of \(\mu_0\) (the droplet \(\Sigma\)), we obtain as a (at least formal) limit as \(n \to +\infty\) of (5.3) an equation of the form (5.1) (resp. (5.2) for \(d = 1\)). Figure 5.1 illustrates this blow up procedure around a point \(x_0\) in \(\Sigma\), the support of \(\mu_0\). The final goal is to derive \(W\) as the governing interaction for the limiting infinite point configurations.
We may observe that although we do have controls on quantities
\[ \left| \nabla h'_{n,\eta} \right|^2 \]
which in turn give controls on the gradient \( \nabla h'_{n,\eta} \) and \( \nabla h'_{n} \) (as seen in Lemma 4.9), we do not fully control \( h'_n \) itself, and thus we will not know its limit \( h \) itself. Also note that (5.1) determines \( h \) from the data of the points only up to a harmonic function.

Since we will work a lot with the gradient of \( h \), it is sometimes convenient to denote it by \( E \), standing for “electric field.” Indeed, \( \nabla h \) physically corresponds to the electric field generated by the charge distribution \( \sum_p N_p \delta_p - m \). The equation (5.1) can then be rewritten with left-hand side \(- \text{div} E\), since \( \text{div} \nabla = \Delta \).

We will give two definitions of the renormalized energy. They rely on the two different ways of subtracting off the self-interaction energy of each charge, that we have already encountered. One definition was first introduced in \[SS6\] for the study of vortices in the Ginzburg-Landau, it relies on the “precursor” \( W \) of Definition 3.1 à la Bethuel-Brezis-Hélein \[BBH\], and was only originally written down in dimensions 1 and 2. The other one was later introduced in \[RouSe\], it relies on the method of smearing out the charges and works in any dimension \( d \geq 2 \) (and certainly \( d = 1 \) as well, although it was not written down).

In the sequel, \( K_R \) will denote the \( d \)-dimensional cubes \([-R,R]^d\).

**Definition 5.1 (Admissible electric fields).** Let \( m \) be a positive number. If \( d \geq 2 \), we let \( \bar{A}_m \) be the set of gradient vector-fields \( E = \nabla h \) that belong to \( L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d) \) for all \( q < \frac{d}{d-1} \), and such that
\[
- \text{div} E = c_d \left( \sum_{p \in \Lambda} N_p \delta_p - m \right) \text{ in } \mathbb{R}^d
\]
for some discrete set \( \Lambda \in \mathbb{R}^d \), and \( N_p \) positive integers; resp. if \( d = 1 \), the set of gradient vector-fields \( E = \nabla h \in L^q_{\text{loc}}(\mathbb{R}^2, \mathbb{R}^2) \) for all \( q < 2 \) such that
\[
- \text{div} E = c_d \left( \sum_{p \in \Lambda} N_p \delta_p - m \delta_R \right) \text{ in } \mathbb{R}^2.
\]
We also let \( A_m \) be the set of gradient vector-fields satisfying the same conditions but with all coefficients \( N_p = 1 \) (i.e. there are no multiple points).

**Definition 5.2 (Renormalized energy by smearing out the charges \[RouSe\]).** Let \( d \geq 2 \). For any \( E \in \bar{A}_m \), and any \( \eta > 0 \) we define
\[
E_\eta = E + \sum_{p \in \Lambda} N_p \nabla f_\eta (\cdot - p)
\]
\[\text{1}\](this is simply the best integrability of the gradient of the Coulomb kernel, as seen in Remark 4.10)
where $\Lambda$ and $N_p$ are associated to $E$ via (5.4), and $f_\eta$ is as in Corollary 4.1; and we let

$$W_\eta(E) = \limsup_{R \to +\infty} \int_{K_R} |E_\eta|^2 - m(\kappa_{d2}(\eta) + \gamma_2 1_{d=2}).$$

(5.6)

We then define the renormalized energy $W$ by

$$W(E) := \liminf_{\eta \to 0} W_\eta(E).$$

(5.7)

We note that one may equivalently define $h_\eta$ to be the smeared out version of the electrostatic potential, that is

$$\nabla h_\eta = \nabla h + \sum_{p \in \Lambda} N_p \nabla f_\eta(\cdot - p),$$

(5.8)

which solves

$$-\Delta h_\eta = c_d \left( \sum_{p \in \Lambda} N_p \delta_p^{(\eta)} - m \right),$$

(5.9)

and then $E_\eta = \nabla h_\eta$.

**Definition 5.3** (Renormalized energy by cutting out holes [SS6,SS8]). Let $d \geq 2$. For any element $E$ of $A_m$, we define the renormalized energy $W$ by

$$W(E) = \limsup_{R \to +\infty} \frac{W(E, \chi_R)}{|K_R|}$$

(5.10)

where, as in (3.10),

$$W(E, \chi_R) = \lim_{\eta \to 0} \int_{\mathbb{R}^2 \setminus \bigcup_{p \in \Lambda} B(p, \eta)} \chi_R |E|^2 - c_d g(\eta) \sum_{p \in \Lambda} \chi_R(p)$$

(5.11)

and $\{\chi_R\}_{R>0}$ is any family of cutoff functions on $\mathbb{R}^2$ such that

$$\chi_R \equiv 1$$

on $K_{R-1}$, $\chi_R \equiv 0$ in $K_R^c$ and $\|\nabla \chi_R\|_{\infty}$ is uniformly bounded. (5.12)

For $d = 1$, we define $W(E)$ in the same way, except $K_R = [-R, R]$, $\{\chi_R\}$ is a family of functions depending only on the first coordinate in $\mathbb{R}^2$, such that $\chi_R \equiv 1$ in $K_{R-1} \times \mathbb{R}$ and $\chi_R \equiv 0$ on $(K_R \times \mathbb{R})^c$, and

$$W(E, \chi_R) = \lim_{\eta \to 0} \int_{\mathbb{R}^2 \setminus \bigcup_{p \in \Lambda} B(p, \eta)} \chi_R |E|^2 + 2\pi \log \eta \sum_{p \in \Lambda} \chi_R(p)$$

(5.13)

Here one should not confuse the precursor to the renormalized energy $W(\cdot, \cdot)$ as in (3.10), which is a function of two entries, with the renormalized energy $W(\cdot)$ that we just defined, and which is a function of one entry.

\footnote{the constants involved were defined in 4.1.5}
We wrote down here an equivalent for dimension \( d \geq 3 \) of the \( W \) defined in [SS6] in dimension 2. It sufficed to replace \( 2\pi \) by \( c_d \) and \( \log \eta \) by \(-g(\eta)\). However, the good properties we need for this energy, such as the fact that it is bounded below and its minimum is achieved, have not been written down anywhere. It is likely that the methods we present here for showing this for \( W \) extend to \( W \) at least up to dimension 2. However, it is not completely clear that we would be able to derive \( W \) from the Coulomb gas Hamiltonian \( H_\alpha \) in dimension \( d \geq 3 \).

### 5.2 First properties

Let us now make a few remarks on the definitions and the comparison between them.

- Both definitions correspond to computing an average energy per unit volume in the limit of a large box size. It is necessary to do so because the system is infinite and otherwise would have an infinite energy.
- It is not a priori clear how to define a total Coulomb interaction of such a “jellium system”, first because of the infinite size of the system as we just saw, second because of the lack of local charge neutrality of the system. The definitions we presented avoid having to go through computing the sum of pairwise interaction between particles (it would not even be clear how to sum them), but instead replace it with (renormalized variants of) the extensive quantity \( \int |\nabla h|^2 \) (see (3.8) and the comments following it).
- In the definition of \( W \), the need for the cut-off functions \( \chi_R \) is due to the fact that if \( \partial K_R \) intersects some ball \( B(p, \eta) \), then the value of \( W \) oscillates wildly between \(+\infty\) and \(-\infty\). Cutting off by a \( C^1 \) function removes the contribution of the points near the boundary, which is generally negligible compared to the total volume anyway.
- The two definitions correspond to two different ways of “renormalizing” i.e. subtracting off the infinite contribution of the Dirac masses to the energy, but more importantly they correspond to reversing the order of the limits \( \eta \to 0 \) and \( R \to +\infty \). As a result the values of \( W \) and \( W \) may differ. This is already seen in the fact that \( W \) accepts multiple points (i.e. \( N_p > 1 \)), while \( W \) is (formally) infinite for multiple points. Indeed, let us suppose that \( N_p \geq 2 \) for some point \( p \in \Lambda \) at distance at least \( 2\eta_0 \) from its neighbors. Then

\[
\int_{B(p, \eta_0)} |E_\eta|^2 \approx \int_{B(0, \eta_0)} |N_p \nabla (g * \delta_0^{(\eta)})|^2 \\
\approx c_d N_p^2 D(\delta_0^{(\eta)}, \delta_0^{(\eta)}) = N_p^2 (\kappa_d g(\eta) + \gamma_2 1_{d=2}).
\]

A few such multiple points add in \( W \) an extra contribution of \( (N_p^2 - N_p)(\frac{\kappa_d}{c_d} g(\eta) + \gamma_2 1_{d=2}) \), which disappears as \( R \to +\infty \) when dividing by \( |K_R| \). On the other
hand, in the second definition,
\[ \int_{B(p,r_0) \setminus B(p,\eta)} |E|^2 \approx N_p^2 c_d g(\eta) \]
which gives an extra contribution to \( W \) of order \( (N_p^2 - N_p)c_d g(\eta) \) and this term diverges to \( +\infty \) as \( \eta \) tends to 0.

- However, as we will see below, \( W \) and \( \hat{W} \) agree for configurations of points which are “well-separated” i.e. for which all the points are simple and separated by a fixed minimum distance, because in that case the order of the limits can be reversed. In addition, in such a situation the value of \( W \) is also independent on the particular choice of function \( \rho \) made for the smearing. We will see below that when dealing with minimizers of these energies, we can reduce to such well-separated configurations.

- The functions \( W \) and \( \hat{W} \) are functions of \( \nabla h \) and not only of the points (recall that \( h \) may vary by addition of a harmonic function), however one can make them functions of the points only by setting, if \( \nu = \sum_{p \in \Lambda} N_p \delta_p \),
\[ \mathcal{W}(\nu) = \inf \{ W(E), -\text{div } E = c_d (\nu - m) \} \]
and the same for \( \hat{W} \), that is taking the infimum of \( W(E) \) on the set of gradient vector-fields \( E \) that are compatible with \( \nu \). Fortunately, these are still measurable as functions of the points, thanks to a measurable selection theorem. For more details, we refer to \[SS7, \text{Sec. 6.6}\].

We next gather some properties of \( W \) and \( \hat{W} \) whose proofs can be found in \[SS6, \text{RouSe}\] respectively.

**Proposition 5.4 (Properties of \( W \) and \( \hat{W} \)).**

1. For \( d = 1, 2 \), the value of \( W \) does not depend on the choice of the family \( \{ \chi_R \}_{R > 0} \) as long as it satisfies (5.12).
2. Both \( W \) and \( \hat{W} \) are Borel-measurable on \( A_m \) (resp. \( \bar{A}_m \)) (and over \( L^q_{\text{loc}} \) when extending the functions by \( +\infty \) outside their domains of definition).
3. Scaling property : if \( E \) belongs to \( A_m \) (resp. \( \bar{A}_m \)), then
\[ \hat{E} := m^{1/d-1} E \left( \frac{\cdot}{m^{1/d}} \right) \in A_1, \text{ resp. } \bar{A}_1. \]

Moreover, we have, if \( d \geq 3 \),
\[
\begin{align*}
W(E) &= m^{2-2/d} W(\hat{E}) \\
W_\eta(E) &= m^{2-2/d} W_{q^{1/a}}(\hat{E})
\end{align*}
\]
and if \( d = 1, 2 \),
\[
\begin{align*}
W(E) &= m W(\hat{E}) - \frac{2\pi}{d} m \log m \\
W(E) &= m W(\hat{E}) - \frac{2\pi}{d} m \log m \\
W_\eta(E) &= m W_{q^{1/a}}(\hat{E}) - \frac{2\pi}{d} m \log m
\end{align*}
\]
One may thus reduce to studying \( W \) and \( \bar{W} \) on \( \mathcal{A}_1 \), resp. \( \bar{\mathcal{A}}_1 \).

4. \( \min_{\bar{\mathcal{A}}_1} W \) is finite and achieved for any \( d \geq 2 \), and \( \min_{\mathcal{A}_1} W \) is finite and achieved for \( d = 1, 2 \). Moreover when \( d = 2 \), the values of these two minima coincide.

5. The minimum of \( W \) on \( \bar{\mathcal{A}}_1 \), resp. of \( W \) on \( \mathcal{A}_1 \) for \( d = 1, 2 \), coincides with the limit as \( N \to +\infty \) of the minima of \( W \) on vector-fields that are \((NZ)^d\)-periodic (i.e. that live on the torus \( \mathbb{T}_N = \mathbb{R}^d/(NZ)^d \)).

It can be expected that in order to balance charges, the constant \( m \), which is the density of the neutralizing background, is also the density of points associated to an \( E \in \bar{\mathcal{A}}_m \). This is in fact true on average for configurations with finite energy. Let us show it in the case of \( W \) which is easier.

**Lemma 5.5.** Let \( E \in \bar{\mathcal{A}}_m \) be such that \( W(E) < +\infty \). Then, letting \( \nu = -\text{div} E + m = \sum_{p \in \Lambda} N_p \delta_p \), we have

\[
\lim_{R \to +\infty} \frac{\nu(K_R)}{|K_R|} = m.
\]

**Proof.** First we show that

\[
\begin{cases}
\nu(K_{R-2}) \leq m|K_R| + CR^{\frac{d-1}{2}}\|E_\eta\|_{L^2(K_R)} \\
\nu(K_{R+1}) \geq m|K_{R-1}| - CR^{\frac{d-1}{2}}\|E_\eta\|_{L^2(K_R)}.
\end{cases}
\] (5.16)

Indeed, by a mean value argument, we may find \( t \in [R-1, R] \) such that

\[
\int_{\partial K_t} |E_\eta|^2 \leq \int_{K_R} |E_\eta|^2.
\] (5.17)

Let us next integrate (5.9) over \( K_t \) and use Green’s formula to find

\[
\int_{K_t} \sum_{p \in \Lambda} N_p \delta_p^{(\eta)} - m|K_t| = - \int_{\partial K_t} E_\eta \cdot \nu,
\] (5.18)

where \( \nu \) denotes the outer unit normal. Using the Cauchy-Schwarz inequality and (5.17), we deduce that

\[
\left| \int_{K_t} \sum_{p \in \Lambda} N_p \delta_p^{(\eta)} - m|K_t| \right| \leq CR^{\frac{d-1}{2}}\|E_\eta\|_{L^2(K_R)}.
\] (5.19)

Since \( \eta \leq 1 \), by definition of \( \nu \) and since the \( \delta_p^{(\eta)} \)'s are supported in the \( B(p, \eta) \)'s, we have \( \nu(K_{R-2}) \leq \int_{K_t} \sum_{p \in \Lambda} N_p \delta_p^{(\eta)} \leq \nu(K_{R+1}) \). The claim (5.16) follows.

Since \( W(E) < +\infty \) then, by definition of \( W \), we have \( W_\eta(E) < +\infty \) for some \( \eta < 1 \) (this is all we really use) and it follows that \( \int_{K_R} |E_\eta|^2 \leq C_\eta R^d \) for any \( R > 1 \). Inserting this into (5.16), dividing by \( |K_R| \) and letting \( R \to \infty \), we easily get the result. \( \square \)
5.3 Well separated configurations and periodic configurations

In this section, we are going to see how to compute explicitly the renormalized energies for periodic configurations of points. In order to do so, we first prove the equivalence of the two ways of computing the renormalized energy for configurations of points which are well-separated.

We start with the following, where the assumptions are the same as in Lemma 3.6.

Lemma 5.6 (The energy of well-separated configurations).

Let \( d \geq 2 \). Assume that \( h \) solves a relation of the form

\[
- \Delta h = c_d \left( \sum_{p \in \Lambda} \delta_p - \mu(x) \right)
\]

in the sense of distributions, for some discrete set \( \Lambda \), and \( \mu \in L^\infty(\Omega) \). Assume that the points are well-separated in the sense that for some \( r_0 > 0 \),

\[
\min \left( \min_{p \neq p'} \min_{p \in \Lambda} \left| p - p' \right|, \min_{p \in \Lambda} \text{dist}(p, \partial \Omega) \right) \geq 2r_0 > 0.
\]

Then, letting \( h_\eta = h + \sum_{p \in \Lambda} f_\eta(\cdot - p) \), we have

\[
\int_\Omega |\nabla h_\eta|^2 - \#(\Lambda \cap \Omega)(\kappa_d g(\eta) + \gamma_2 1_{d=2}) = W(\nabla h, 1_\Omega) + \#(\Lambda \cap \Omega) o_\eta(1) ||\mu||_{L^\infty(\Omega)},
\]

where \( o_\eta(1) \to 0 \) as \( \eta \to 0 \) is a function that depends only on the dimension.

Remark 5.7. An inspection of the proof shows that the result is true with appropriate modification for \( \mu \) as in Remark 3.3.

Proof. Since \( h_\eta \) is defined by (5.8), the \( B(p, r_0) \) are disjoint and included in \( \Omega \), and \( f_\eta \) is identically 0 outside of \( B(0, \eta) \) we may write for any \( \eta < r_0 \), and any \( 0 < \alpha < \eta \),

\[
\int_{\Omega \setminus \bigcup_{p \in \Lambda} B(p, \alpha)} |\nabla h_\eta|^2 = \int_{\Omega \setminus \bigcup_{p \in \Lambda} B(p, \alpha)} |\nabla h|^2 + \#(\Lambda \cap \Omega) \int_{B(0, \alpha) \setminus B(0, \alpha)} |\nabla f_\eta|^2 + 2 \sum_{p \in \Lambda} \int_{B(p, \eta) \setminus B(p, \alpha)} \nabla f_\eta(x - p) \cdot \nabla h. \quad (5.23)
\]

First we note that, using Green’s formula, and \( \nu \) denoting the outwards pointing unit normal to \( \partial B(0, \alpha) \), we have

\[
\int_{B(0, \eta) \setminus B(0, \alpha)} |\nabla f_\eta|^2 = - \int_{\partial B(0, \alpha)} f_\eta \frac{\partial f_\eta}{\partial \nu} + c_d \int_{B(0, \alpha) \setminus B(0, \alpha)} f_\eta g_\eta^{(\alpha)}.
\]
By Green’s formula again and the definition of \( f_\eta \) we have
\[
\int_{\partial B(0,\alpha)} \frac{\partial f_\eta}{\partial \nu} = -c_d \int_{B(0,\alpha)} \delta_0^{(\eta)} + c_d = c_d + o_\alpha(1) \quad \text{as } \alpha \to 0
\]
and combining with the fact that \( f_\eta = g * \delta_0^{(\eta)} - g \), we find
\[
\int_{B(0,\eta) \setminus B(0,\alpha)} \left| \nabla f_\eta \right|^2 = -c_d (g * \delta_0^{(\eta)}(\alpha) - g(\alpha)) + c_d \int_{\mathbb{R}^d} \left( g * \delta_0^{(\eta)} \right) \delta_0^{(\eta)} - c_d \int_{\mathbb{R}^d} g \delta_0^{(\eta)} + o_\alpha(1). \quad (5.24)
\]
We next observe that \( c_d \int_{\mathbb{R}^d} \left( g * \delta_0^{(\eta)} \right) \delta_0^{(\eta)} = c_d D(\delta_0^{(\eta)}, \delta_0^{(\eta)}) \) and \( \int_{\mathbb{R}^d} g \delta_0^{(\eta)} = g * \delta_0^{(\eta)}(0) \), thus, inserting into (5.24) and using (4.16), we find
\[
\int_{B(0,\eta) \setminus B(0,\alpha)} \left| \nabla f_\eta \right|^2 = -2c_d g * \delta_0^{(\eta)}(0) + c_d g(\alpha) + \kappa_d g(\eta) + \gamma_2 1_{d=2} + o_\alpha(1),
\]
in view of the fact that for fixed \( \eta \), \( g * \delta_0^{(\eta)} \) is continuous at 0. On the other hand, using Green’s formula and (5.20) we have
\[
\int_{B(\eta) \setminus B(p,\alpha)} \nabla f_\eta(x-p) \cdot \nabla h = -c_d \int_{B(\eta) \setminus B(p,\alpha)} f_\eta(x-p) \mu(x) \, dx - f_\eta(\alpha) \int_{\partial B(p,\alpha)} \frac{\partial h}{\partial \nu}.
\]
First we note that
\[
\left| \int_{B(\eta) \setminus B(p,\alpha)} f_\eta(x-p) \mu(x) \, dx \right| \leq \| \mu \|_{L^\infty} o_\alpha(1),
\]
where \( o_\alpha(1) \) depends only on \( d \). To see this, just notice that we have \( |f_\eta| \leq |g| \) and the Coulomb kernel \( g \) is integrable near the origin. Secondly, by Green’s theorem again, we have
\[
-\int_{\partial B(p,\alpha)} \frac{\partial h}{\partial \nu} = c_d + O(\| \mu \|_{L^\infty} \alpha^d).
\]
Inserting these two facts we deduce
\[
\int_{B(\eta) \setminus B(p,\alpha)} \nabla f_\eta(x-p) \cdot \nabla h = c_d f_\eta(\alpha) + O(\| \mu \|_{L^\infty} \alpha^d g(\alpha)) + \| \mu \|_{L^\infty} o_\alpha(1). \]
Combining this and (5.25), (5.23) and again \( f_\eta(\alpha) = g * \delta_0^{(\eta)}(0) - g(\alpha) + o_\alpha(1) \), we find
\[
\int_{B(\eta) \setminus \bigcup_{\lambda \in \Lambda} B(p,\alpha)} |\nabla h_\eta|^2 = \int_{B(\eta) \setminus \bigcup_{\lambda \in \Lambda} B(p,\alpha)} |\nabla h|^2 + \#(\Lambda \cap \Omega) \left[ \kappa_d g(\eta) + \gamma_2 1_{d=2} - c_d g(\alpha) + o_\alpha(1) + \| \mu \|_{L^\infty} o_\eta(1) + \| \mu \|_{L^\infty} O(\alpha^d g(\alpha)) \right].
\]
5.3 Well separated configurations and periodic configurations

Letting $\alpha \to 0$, in view of the definition of $W$ in (3.10), we obtain the result.

\[\]

\textbf{Corollary 5.8} ($W$ and $\mathcal{W}$ coincide in 2D for well-separated points).

Assume $d = 2$, and let $E \in \mathcal{A}_1$ be such that $\mathcal{W}(E) < +\infty$ and the associated set of points satisfies $\min_{p \neq p' \in \Lambda} |p - p'| \geq 2r_0 > 0$ for some $r_0 > 0$. Then $\mathcal{W}(E) = W(E)$.

For the proof, see [RouSe, Prop. 3.3]. It is very likely that this can be extended at least to dimension $d = 3$.

We now turn to periodic configurations, and show that, for them, $W$ or $\mathcal{W}$ can be computed and expressed as a sum of pairwise Coulomb-like interactions between the points. By periodic configuration, we mean a configuration on the fundamental cell of a torus, repeated periodically, which can be viewed as a configuration of $N$ points on a torus (cf. Fig. 5.2).

![Figure 5.2. Periodic configurations](image)

**Proposition 5.9.** Let $a_1, \ldots, a_N$ be $N$ points in a torus $\mathbb{T}$ of volume $N$ in $\mathbb{R}^d$, $d \geq 2$.

1. If there is a multiple point, then for any $E$ compatible with the points (i.e. such that $-\text{div } E = c_d (\sum_{i=1}^N a_i - 1)$), we have $W(E) = +\infty$.

2. If all points are distinct, letting $H$ be the periodic solution to

\[ -\Delta H = c_d \left( \sum_{i=1}^N \delta_{a_i} - 1 \right), \quad \int_\mathbb{T} H = 0; \tag{5.26} \]

then any other periodic $E$ compatible with the points satisfies

\[ W(E) \geq W(\nabla H). \tag{5.27} \]
Moreover,
\[
\mathcal{W}(\nabla H) = \frac{c_d^2}{N} \sum_{i \neq j} G(a_i - a_j) + c_d^2 \lim_{x \to 0} \left( G - \frac{g}{c_d} \right) \tag{5.28}
\]
where \( G \), the Green function of the torus, solves
\[
-\Delta G = \delta_0 - \frac{1}{|T|} \text{ over } \mathbb{T}, \quad \int_T G = 0. \tag{5.29}
\]

3. If \( d = 2 \), the same results hold true with \( W \) instead of \( \mathcal{W} \).

4. If \( d = 1 \), the same results hold true with \( W \), \( c_d = 2\pi \) and \( H \) the solution on \( \mathbb{R}/(N\mathbb{Z}) \times \mathbb{R} \) with zero-mean on the real axis of
\[
-\Delta H = 2\pi \left( \sum_{i=1}^{N} \delta_{a_i} - \delta_R \right),
\]
\[
G(x) = -\frac{1}{2\pi} \sum_{i \neq j} \log \left| 2\sin \frac{\pi x}{N} \right|, \quad g(x) = -\log |x|, \tag{5.30}
\]
i.e. we have
\[
\mathcal{W}(\nabla H) = -\frac{2\pi}{N} \sum_{i \neq j} \log \left| 2\sin \frac{\pi(a_i - a_j)}{N} \right| - 2\pi \log \frac{2\pi}{N}. \tag{5.31}
\]

Proof. We start with the first assertion: let \( N_i \) be the multiplicity of \( a_i \). We have \( \sum_i N_i = N \). On the torus, for any \( E \) compatible with the points, we may check using similar calculations as in the proof of Lemma 5.6 that
\[
\int_T |E_\eta|^2 \geq \sum_i N_i^2 \kappa_d g(\eta) - CN.
\]
Meanwhile, by periodicity, we have that
\[
\mathcal{W}(E) = \lim_{\eta \to 0} \left( \frac{1}{|T|} \int_T |E_\eta|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \right).
\]
Since the volume \( |T| = N \), it follows that
\[
\mathcal{W}(E) \geq \frac{1}{N} \left( \sum_i N_i^2 - N \right) \kappa_d g(\eta) - C
\]
and the limit as \( \eta \to 0 \) of this quantity is \(+\infty\) unless \( \sum_i N_i^2 = N \), which imposes that all the multiplicities \( N_i \) be equal to 1. The same goes for \( W \) (for which we already know that multiple points give infinite value).
5.3 Well separated configurations and periodic configurations

For the second assertion, if all the points are simple and the configuration is periodic, it follows that the points are well-separated, i.e. satisfy the assumptions of Lemma 5.6. Thus we know that

$$\int_{\mathbb{T}} |E_\eta|^2 - N(c_d g(\eta) + \gamma_2 1_{d=2}) = W(E, 1_\mathbb{T}) + o_\eta(1) \quad (5.32)$$

and by definition of $W$ and $W$ this immediately proves the identity

$$W(E) = \frac{W(E, 1_\mathbb{T})}{|\mathbb{T}|}. \quad (5.33)$$

At this point, one can also check that in dimension $d = 2$, $W(E)$ is also equal to $W(E, 1_\mathbb{T})$ (this requires a little more care to show that the effect of the cut-off function is negligible, see [SS6] for details).

Assume now that $E_1$ and $E_2$ are two admissible periodic gradient vector-fields, with $E_1 = \nabla h_1$ and $E_2 = \nabla h_2$. Since $\Delta(h_1 - h_2) = 0$, $E_1$ and $E_2$ differ by the gradient of a harmonic function, but they are also periodic so this difference must in fact be a constant vector $\vec{c}$. We can then compute

$$W(E_1, 1_\mathbb{T}) = \lim_{\eta \to 0} \int_{\mathbb{T} \cup \partial B(a, \eta)} |E_1 + \vec{c}|^2 - c_d N g(\eta)$$

$$= W(E_2, 1_\mathbb{T}) + \lim_{\eta \to 0} \left[ \int_{\mathbb{T} \cup \partial B(a, \eta)} |\vec{c}|^2 + 2\vec{c} \cdot \int_{\mathbb{T} \cup \partial B(a, \eta)} E_2 \right]$$

$$= W(E_2, 1_\mathbb{T}) + \int_{\mathbb{T}} |\vec{c}|^2 + 2\vec{c} \cdot \int_{\mathbb{T}} E_2. \quad (5.34)$$

If $E_2 = \nabla H$ for some $H$ periodic, then $\int_{\mathbb{T}} E_2 = \int_{\mathbb{T}} \nabla H = 0$, hence we find $W(E_1, 1_\mathbb{T}) \geq W(E_2, 1_\mathbb{T})$. In view of (5.33), this proves (5.27).

Let us now turn to the proof of (5.28). Let $H$ be the periodic solution with mean zero. It is easy to see that $H(x) = c_d \sum_{i=1}^{N} G(x - a_i)$ with $G$ the Green function defined in the proposition. One can then compute the energy in terms of $G$, recalling that

$$W(\nabla H, 1_\mathbb{T}) = \lim_{\eta \to 0} \left( \int_{\mathbb{T} \cup \partial B(a, \eta)} |\nabla H|^2 - N c_d g(\eta) \right)$$

$$= \lim_{\eta \to 0} \left( \sum_{i=1}^{N} \int_{\partial B(a, \eta)} \frac{\partial H}{\partial \nu} - c_d \int_{\mathbb{T} \cup \partial B(a, \eta)} H - N c_d g(\eta) \right). \quad (5.35)$$

The term $\int_{\mathbb{T} \cup \partial B(a, \eta)} H$ converges to $\int_{\mathbb{T}} H = 0$ as $\eta$ goes to zero. Plugging in $H(x) = c_d \sum_{i=1}^{N} G(x - a_i)$ gives that

$$\int_{\partial B(a, \eta)} \frac{\partial H}{\partial \nu} = c_d \sum_{j \neq i} G(a_j - a_i) \int_{\partial B(a, \eta)} \frac{\partial H}{\partial \nu} + c_d G(\eta) \int_{\partial B(a, \eta)} \frac{\partial H}{\partial \nu} + o_\eta(1)$$
and an application of Green’s formula allows us to rewrite $\int_{\partial B(a_i, \eta)} \frac{\partial H}{\partial \nu}$ as $-\int_{B(a_i, \eta)} \Delta H$
and this term equals $c_d + o_\eta(1)$. This gives (with $|T| = N$):

$$W(\nabla H) = \frac{W(\nabla H, 1_T)}{|T|} = \frac{c_d^2}{N} \sum_{i \neq j} G(a_i - a_j) + c_d^2 \left( \frac{G(\eta) - g(\eta)}{c_d} \right) + o_\eta(1)$$

and letting $\eta \to 0$ gives (5.28). The proofs of items 3 and 4 are very similar, we refer to [SS6, SS8] for details. In the case of dimension 1, the suitable Green function $G$ can be explicitly solved by Fourier series, and one finds the formula (5.30).

Note that, in dimension $d \geq 2$, equation (5.29) can also be solved somewhat explicitly. For a torus $T = \mathbb{R}^d / (\mathbb{Z} u_1 + \cdots + \mathbb{Z} u_d)$ of volume $N$, corresponding to the lattice $\Lambda = \mathbb{Z} u_1 + \cdots + \mathbb{Z} u_d$ in $\mathbb{R}^d$, one may first express $G$ solving

$$-\Delta G = \delta_0 - \frac{1}{|T|} = \delta_0 - \frac{1}{N}, \quad \int_T G = 0$$

as a Fourier series :

$$G = \sum_{\vec{k} \in \Lambda^*} c_{\vec{k}} e^{2\pi i \vec{k} \cdot \vec{x}}$$

where $\Lambda^*$ is the dual lattice of $\Lambda$, that is

$$\Lambda^* = \{ \vec{q} \in \mathbb{R}^d, \vec{q} \cdot \vec{p} \in \mathbb{Z} \text{ for all } \vec{p} \in \Lambda \}.$$ 

Plugging this into (5.36), one sees that the coefficients $c_{\vec{k}}$ must satisfy the relations

$$-(2\pi)^2 |\vec{k}|^2 c_{\vec{k}} = 1 - \delta_{\vec{k},0}$$

where $\delta_{\vec{k},0}$ is 0 unless $\vec{k} = 0$, and $c_0 = \int_T G = 0$ by assumption. This is easily solved by $c_{\vec{k}} = \frac{1}{4\pi^2 |\vec{k}|^2}$ for $\vec{k} \neq 0$, hence the formula

$$G(\vec{x}) = \sum_{\vec{k} \in \Lambda^* \setminus \{0\}} \frac{e^{2\pi i \vec{k} \cdot \vec{x}}}{4\pi^2 |\vec{k}|^2}.$$ 

Such a series is called an Eisenstein series, cf. [Lan] for reference and formulas on Eisenstein series.

5.4 Partial results on the minimization of $W$ and $W$

We have seen in item 5 of Proposition 5.4 that the minima of $W$ and $W$ can be achieved as limits of the minima over periodic configurations (with respect to
large and larger tori). On the other hand, Proposition 5.9 provides a more explicit expression for periodic configurations. In dimension \( d = 1 \) (and in that case only) we know how to use this expression (5.31) to identify the minimum over periodic configurations: a convexity argument (for which we refer to [SS8, Prop. 2.3]) shows that the minimum is achieved when the points are equally spaced, in other words for the lattice or crystalline distribution \( \mathbb{Z} \) (called “clock distribution” in the context of orthogonal polynomials, cf. [Sim]). Combining with the result of item 5 of Proposition 5.4 allows to identify \( \min_{A_1} W \):

**Theorem 5.1** (The regular lattice is the minimizer in 1D [SS8].) If \( d = 1 \), we have

\[
\min_{A_1} W = -2\pi \log(2\pi)
\]

and this minimum is achieved by gradients of periodic potentials \( h \) associated to the lattice (or clock) distribution \( \Lambda = \mathbb{Z} \).

Of course, the minimum over any \( A_m \) is deduced from this by scaling (cf. (5.15)).

In higher dimension, determining the value of \( \min W \) or \( \min \mathcal{W} \) is an open question, even though it would suffice to be able to minimize in the class of periodic configurations with larger and larger period, using the formula (5.28). The only question that we can answer so far is that of the minimization over the restricted class of pure lattice configurations, in dimension \( d = 2 \) only, i.e. vector fields which are gradient of functions that are periodic with respect to a lattice \( \mathbb{Z}\vec{u} + \mathbb{Z}\vec{v} \) with \( \det(\vec{u}, \vec{v}) = 1 \), corresponding to configurations of points that can be identified with \( \mathbb{Z}\vec{u} + \mathbb{Z}\vec{v} \). In this case, we have:

**Theorem 5.2** (The triangular lattice is the minimizer over lattices in 2D [SS6].) The minimum of \( W \), or equivalently \( \mathcal{W} \), over this class of vector fields is achieved uniquely by the one corresponding to the triangular “Abrikosov” lattice.

Here the triangular lattice means \( \mathbb{Z} + \mathbb{Z}e^{i\pi/3} \), properly scaled, i.e. what is called the Abrikosov lattice in the context of superconductivity, cf. Chap. 1.

We next give a sketch of the proof, which is not very difficult thanks to the fact that it reduces to the same question for a certain modular function, which was solved by number theorists in the 50’s and 60’s.

**Proof of Theorem 5.2** Proposition 5.9 more specifically (5.28), provides an explicit formula for the renormalized energy of such periodic configurations. Using (5.38) to express \( G \), and denoting by \( H_\Lambda \) the periodic solution associated with (5.26), we find that

\[
\mathcal{W}(\nabla H_\Lambda) = \lim_{x \to 0} \left( \sum_{\vec{k} \in \Lambda \setminus \{0\}} \frac{e^{2i\pi \vec{k} \cdot \vec{x}}}{4\pi^2 |\vec{k}|^2} + 2\pi \log x \right).
\]
By using either the “first Kronecker limit formula” (cf. [Lan]) or a direct computation, one shows that in fact

$$W(\nabla H_\Lambda) = C_1 + C_2 \lim_{x \to 0, x > 0} \left( \sum_{k \in \Lambda^\ast \setminus \{0\}} \frac{1}{|k|^2 + x} - \int_{\mathbb{R}^2} \frac{dy}{1 + |y|^2 + x} \right),$$

(5.40)

where $C_1$ and $C_2 > 0$ are constants. The series $\sum_{k \in \Lambda^\ast \setminus \{0\}} \frac{1}{|k|^2 + x}$ that appears is now the “Epstein Zeta function” of the dual lattice $\Lambda^\ast$. The first Kronecker limit formula allows to pass from one modular function, the Eisenstein series, to another, the Epstein Zeta function. Note that both formulas (5.39) and (5.40), when $x \to 0$, correspond to two different ways of regularizing the divergent series $\sum_{p \in \Lambda^\ast \setminus \{0\}} \frac{1}{|p|^2}$, and they are in fact explicitly related.

The question of minimizing $W$ among lattices is then reduced to minimizing the Epstein Zeta function

$$\Lambda \mapsto \zeta_\Lambda(x) := \sum_{k \in \Lambda \setminus \{0\}} \frac{1}{|k|^2 + x}$$

as $x \to 0$. But results from [Cas,Ran,Enno1,Enno2,Dia,Mont] assert that

$$\zeta_\Lambda(x) \geq \zeta_{\Lambda_{\text{triang}}}(x), \quad \forall x > 0$$

(5.41)

and the equality holds if and only if $\Lambda = \Lambda_{\text{triang}}$ (the triangular lattice). Because that lattice is self-dual, it follows that it is the unique minimizer.

The same result was obtained in [CO] for a similar energy. In dimension $d \geq 3$ the same computation holds but the meaning of (5.40) is not clear. The minimization of the Epstein Zeta function over lattices is then an open question (except in dimensions 8 and 24). In dimension 3, both the FCC (face centered cubic) and BCC (boundary centered cubic) lattices (cf. Fig. 5.3) could play the role of the triangular lattice, but it is only conjectured that FCC is a local minimizer (cf. [SaSt]).

Figure 5.3. BCC and FCC lattices
One may ask whether this triangular lattice does achieve the global minimum of $W$ and $\mathcal{W}$. The fact that the Abrikosov lattice is observed in superconductors, combined with the fact – which we will see later – that $W$ can be derived as the limiting minimization problem of Ginzburg-Landau, justify to conjecture this:

**Conjecture 5.1.** In dimension $d = 2$, the value of $\min_{A_1} W = \min_{A_1} \mathcal{W}$ is equal to the value at the vector field associated to the triangular lattice of volume 1.

This belongs to the wider class of crystallization problems. A typical question of this sort is, given a potential $V$ in any dimension, to determine the point positions that minimize

$$\sum_{i \neq j} V(x_i - x_j)$$

(with some kind of boundary condition), or rather

$$\lim_{R \to \infty} \frac{1}{|B_R|} \sum_{i \neq j, x_i, x_j \in B_R} V(x_i - x_j),$$

and to determine whether the minimizing configurations are perfect lattices. Such questions are fundamental in order to understand the crystalline structure of matter. They also arise in the arrangement of Fekete points [SK], “Smale’s 7th problem” on the sphere, or the “Cohn-Kumar conjecture” [CoKu]. One should immediately stress that there are very few positive results in that direction in the literature (in fact it is very rare to have a proof that the solution to any minimization problem is periodic). Some exceptions include the two-dimensional sphere packing problem, for which Radin [Ra] showed that the minimizer is the triangular lattice, and an extension of this by Theil [Th] for a class of very short range Lennard-Jones potentials. The techniques used there do not apply to Coulomb interactions, which are much longer range. Let us mention another recent positive result. The question of minimization of $\mathcal{W}$ can also be very informally rephrased as that of finding

$$\min \left\| \sum_p \delta_p - 1 \right\|_{(H^1)'}$$

where the quantity is put between brackets to recall that $\delta_p$ does not really belong to the dual of the Sobolev space $H^1$ but rather has to be computed in the renormalized way that defines $\mathcal{W}$. A closely related problem is to find

$$\min \left\| \sum_p \delta_p - 1 \right\|_{\text{Lip}^*},$$

and it turns out to be much easier. It is shown by Bourne-Peletier-Theil in [BPT] with a relatively short proof that again the triangular lattice achieves the minimum.

We finish by referring to some extra results.
With Rota Nodari, in [RNSe], we showed the equivalence between several ways of phrasing the minimization of \( W \) in dimension 2 over a finite size box: minimization with prescribed boundary trace and minimization among periodic configurations. In all cases, we were able to prove, in the spirit of [ACO], that the energy density and the points were uniformly distributed at any scale \( \gg 1 \), in good agreement with (but of course much weaker than!) the conjecture of periodicity of the minimizers.

Even though the minimization of \( W \) is only conjectural, it is natural to view it as (or expect it to be) a quantitative “measure of disorder” of a configuration of points in the plane. In this spirit, with Borodin [BSe], we used \( W \) (or rather a variant of it) in dimensions 1 and 2 to quantify and compute explicitly the disorder of some classic random point processes in the plane and on the real line.

## 5.5 Lower bound on \( W \)

We saw in item 4 of Proposition 5.4 that \( \min_{\bar{A}_1} W \) is finite. This property, in addition to its intrinsic interest, will turn out to be crucial for us in the next chapter when deriving rigorously \( W \) from the Coulomb gas Hamiltonian, in the limit \( n \to \infty \).

In this section we will give an idea of how to prove a lower bound on \( W \). This is a place where the analysis developed for \( W \) in [RouSe], via the smearing out of charges, differs from that developed for \( W \) in [SS6]: there it is also proved that \( W \) is bounded below, in dimension \( d = 2 \) (it also works for \( d = 1 \)) but by a different method, relying on the “ball construction”, à la Jerrard [Je] and Sandier [Sa], which only works in dimensions 1 and 2. It is likely that the method we present below, which follows [RouSe], can also be adapted for \( W \) in dimensions 2 and 3 at least, but not all details have been checked.

Let us state the main result.

**Proposition 5.10.** Let \( d \geq 2 \). For \( \eta \leq 1 \), the functions \( \mathcal{W}_\eta \) and \( W \) are bounded below on \( \bar{A}_1 \) by the same constant depending only on the dimension \( d \).

The method consists in first reducing to configurations with points that are simple and “well-separated,” in the sense seen previously. This relies on an unpublished result of E. Lieb [Lie2], which can be found in dimension \( d = 2 \) in [RNSe, Theorem 4] and which we readapt to the setting of smeared out charges, i.e. for any \( \eta \leq 1 \). This step is accomplished via the following

**Proposition 5.11 (Reducing to well-separated points).** Let \( \Lambda \) be a discrete subset of \( K_R \) and let \( h \) satisfy

\[
- \Delta h = c_d \left( \sum_{p \in \Lambda} N_p \delta_{\eta} \quad - 1 \right) \quad \text{in} \quad K_R.
\] (5.42)
5.5 Lower bound on $\mathcal{W}$

Denote $\Lambda_R = \Lambda \cap K_{R-1}$. There exist three positive constants $\eta_0, r_0, C$ such that if $\eta < \eta_0$, $R$ is large enough and one of the following conditions does not hold:

\begin{align*}
\forall p \in \Lambda_R, \quad N_p &= 1 \quad (5.43) \\
\forall p \in \Lambda_R, \quad \text{dist}(p, \Lambda_R \setminus \{p\}) &\geq 2r_0, \quad (5.44)
\end{align*}

then there exists $\hat{\Lambda}$, a discrete subset of $K_R$ and an associated potential $\hat{h}$ satisfying

\begin{equation}
-\Delta \hat{h} = cd \left( \sum_{p \in \hat{\Lambda}} N_p \delta^{(\eta)}_p - 1 \right) \quad \text{in } K_R 
\end{equation}

such that

\begin{equation}
\int_{K_R} |\nabla \hat{h}|^2 \leq \int_{K_R} |\nabla h|^2 - C.
\end{equation}

Another way of phrasing the proposition is that if a configuration has points that are too close to one another, we can always replace it by one that has a smaller energy, thus we can always take a minimizing sequence for

\begin{equation}
F_{\eta, R} := \inf \left\{ \int_{K_R} |\nabla h|^2, -\Delta h = cd \left( \sum_{p \in \Lambda} N_p \delta^{(\eta)}_p - 1 \right) \quad \text{in } K_R \right\}.
\end{equation}

with points that are simple and well-separated (at least in $K_{R-1}$). Note that in the proposition, the configuration of points $\Lambda$ may depend on $\eta$.

The second step is what we call the “screening” of a configuration. Once we have reduced to such a configuration with well-separated points, we show that we can efficiently “screen” it. This means modify it near the boundary of the cube $K_R$ to make the normal component of the electric field vanish on $\partial K_R$, still keeping the points well-separated all the way to the boundary of the cube. This modification needs to add only a negligible energy cost. The vanishing normal component will in particular impose the total number of points in the cube, but it also makes the configurations “boundary compatible” with each other, which will allow to copy and paste them together, e.g. to periodize them. Physically, “screening” roughly means here that a particle sitting outside of $K_R$ does not “feel” any electric field coming from $K_R$.

**Proposition 5.12 (Screening).** There exists $\eta_0 > 0$ such that the following holds for all $\eta < \eta_0$. Let $h_\eta$ satisfy $\bbox[1.5pt,red,1pt]{5.42} - \bbox[1.5pt,blue,1pt]{5.43} - 5.44$ and

\begin{equation}
\int_{K_R} |\nabla h_\eta|^2 \leq C_\eta R^d.
\end{equation}

Then there exists $\hat{\Lambda}$ a configuration of points and $\nabla \hat{h}$ an associated gradient vector field (both possibly also depending on $\eta$) defined in $K_R$ and satisfying

\begin{equation}
\begin{cases}
-\Delta \hat{h} = cd \left( \sum_{p \in \hat{\Lambda}} \delta_p - 1 \right) \quad \text{in } K_R \\
\frac{\partial \hat{h}}{\partial \nu} = 0 \quad \text{on } \partial K_R
\end{cases}
\end{equation}
such that for any $p \in \hat{\Lambda}$

$$\min \left( \text{dist}(p, \hat{\Lambda} \setminus \{p\}), \text{dist}(p, \partial K_R) \right) \geq \frac{r_0}{10}$$

(5.49)

with $r_0$ as in (5.44), and

$$\int_{K_R} |\nabla \hat{h}_\eta|^2 \leq \int_{K_R} |\nabla h_\eta|^2 + o(R^d)$$

(5.50)

as $R \to \infty$, where the $o$ depends only on $\eta$.

Remark 5.13. The vanishing of the normal derivative $\partial h/\partial \nu$ on $\partial K_R$ implies, by Green’s theorem, that $\#(\hat{\Lambda} \cap K_R) = |K_R|$ exactly.

Remark 5.14. In [LS], we show that the $\liminf_{\eta \to 0} W_\eta$ that defines $W$ is an increasing limit, modulo a universal $o_\eta(1)$, and we deduce a short proof of Proposition 5.10.

Let us start by giving the idea of the proof of Proposition 5.12. For details, cf. [RouSe].

Proof. Step 1. It consists, by a mean value argument, to select a “good boundary” $\partial K_t$ (cf. Fig. 5.4) at distance $L$ with $1 \ll L \ll R$ (as $R \to \infty$) from $\partial K_R$, and such that

$$\left\{ \begin{array}{l}
\int_{\partial K_t} |\nabla h_\eta|^2 \leq \frac{C_\eta R^d}{L} = o(R^d) \\
\int_{K_{t+1} \setminus K_{t-1}} |\nabla h_\eta|^2 \leq \frac{C_\eta R^d}{L} = o(R^d).
\end{array} \right.$$  

(5.51)

Step 2. Taking $\eta < r_0/4$ so that the balls $B(p, \eta)$ are disjoint, we can modify the boundary $\partial K_t$ into $\partial \Gamma$ (cf. Fig. 5.4) so that $\partial \Gamma$ intersects no ball $B(p, \eta)$ and still satisfies

$$\int_{\partial \Gamma} |\nabla h_\eta|^2 \leq o(R^d).$$

(5.52)

We do not move the points whose associated smeared charges intersect $\partial K_t$. Instead, we isolate them in small cubes and leave unchanged all the points lying in $\Gamma$, defined as the union of $K_t$ with these small cubes.

Step 3. We build a new configuration of points and potential in $K_R \setminus \Gamma$, to replace the previous one. To do so, we partition this region into hyperrectangles $K_i$, each centered at some point $x_i$, on each of which we solve

$$\left\{ \begin{array}{l}
-\Delta u_i = c_d(\delta_{x_i} - 1) \text{ in } K_i \\
\text{the normal derivatives } \partial u_i/\partial \nu \text{ are compatible}
\end{array} \right.$$  

(5.53)

so that the normal derivatives “connect” nicely, meaning that they agree (with suitable orientation) on any two adjacent hyperrectangles, as well as on the boundary of $\Gamma$. The new set $\hat{\Lambda}$ is defined as $\cup_i \{x_i\} \cup (\Lambda \cap K_{R-1})$. One checks that the
hyperrectangles have sidelengths which are bounded below in such a way that the new set $\hat{\Lambda}$ satisfies (5.49)

**Step 4.** We define a global vector field and estimate its energy. First of all, to $\nabla h_\eta$ given in the statement of the proposition corresponds a $\nabla h$ via (5.8). Since $\Lambda$ may possibly depend on $\eta$, so does $\nabla h$, but this is not important. Defining then $E$ to be $\nabla u_i$ on each $K_i$, and $E = \nabla h$ in $\Gamma$, thanks to the compatibility condition we may check that $E$ satisfies

$$- \text{div } E = c_d \left( \sum_{p \in \hat{\Lambda}} \delta_p - 1 \right) \text{ in } K_R. \quad (5.54)$$

Indeed, one may check that the divergence of $E$ (or any vector field) in the sense of distributions on each interface is given by the jump in normal derivative (here constructed to be zero) while the divergence on each cell $K_i$ is given by (5.53).

By elliptic estimates, and using (5.52), one can evaluate $\int |\nabla u_i|^2$ and we claim that such a construction can be achieved with

$$\int_{K_R \setminus \Gamma} |E_\eta|^2 \leq \sum_i \int_{K_i} |\nabla u_i,\eta|^2 \leq o(R^d),$$

with $E_\eta = E + \sum_{p \in \hat{\Lambda}} \nabla f_\eta(-p)$, i.e. the modification in the boundary layer $K_R \setminus \Gamma$.
can be made with a negligible energy, so that
\[
\int_{K_R} |E_\eta|^2 \leq \int_{K_R} \left| \nabla h_\eta \right|^2 + o(R^d). \tag{5.55}
\]
We would like to define \( \nabla \hat{h} \) as \( E_\eta \), but the problem is that \( E_\eta \) is not a gradient. To remedy this, we use a kind of Hodge (or Helmholtz if \( d \leq 3 \)) decomposition, which consists in adding a vector field to \( E \) to make it a gradient, without changing its divergence, while not deteriorating the energy estimate (5.55). Let us now show this more precisely.

The Hodge decomposition tells us that we may find a vector field \( X \) defined over \( K_R \), such that \( \text{div } X = 0 \) in \( K_R \), the normal component \( X \cdot \nu = 0 \) on \( \partial K_R \) and \( E_\eta + X \) is the gradient of a function, which we call \( \hat{h} \), hence so is \( E_\eta + X = \nabla \hat{h}_\eta \).

An easy computation then yields
\[
\int_{K_R} |E_\eta|^2 = \int_{K_R} |E_\eta + X|^2 - 2 \int_{K_R} (E_\eta + X) \cdot X
\]
But we can apply Green’s theorem on the right-most term and find
\[
\int_{K_R} (E_\eta + X) \cdot X = 0
\]
since we saw that \( E_\eta + X \) is a gradient and by assumption \( \text{div } X = 0 \) and \( X \cdot \nu = 0 \).

We conclude that
\[
\int_{K_R} |E_\eta|^2 \geq \int_{K_R} |E_\eta + X|^2 = \int_{K_R} \left| \nabla \hat{h}_\eta \right|^2 \tag{5.56}
\]
which combined with (5.55), proves that \( \hat{h} \) satisfies all the desired properties (cf. (5.54)).

At this point, we can see how Propositions 5.11 and 5.12 lead to the lower bound stated in Proposition 5.10.

**Corollary 5.15.** Given \( \eta < 1 \), for any \( R \) large enough and such that \( |K_R| \in \mathbb{N} \), there exists an \( \hat{h} \) satisfying (5.48)–(5.49), such that
\[
\limsup_{R \to +\infty} \int_{K_R} \left| \nabla \hat{h}_\eta \right|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq \inf_{\mathcal{A}_1} W_\eta. \tag{5.57}
\]

**Proof.** In view of Propositions 5.11 and 5.12 for each given \( R \), we may choose a \( \nabla \hat{h} \) approximating \( F_{\eta,R} \) and associated to simple well-separated points (separated by \( r_0 \) which may depend only on \( d \)) and with \( \frac{\partial h}{\partial n} = 0 \) on \( \partial K_R \), which implies by Remark 5.13 that \( \#(\Lambda \cap K_R) = |K_R| \). Then, letting \( R \to \infty \), we have
\[
\limsup_{R \to +\infty} \int_{K_R} \left| \nabla \hat{h}_\eta \right|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq \limsup_{R \to +\infty} \frac{F_{\eta,R}}{|K_R|} - (\kappa_d g(\eta) + \gamma_2 1_{d=2}).
\]
Then (5.57) follows as a consequence of the definitions of \( W_\eta \) and \( F_{\eta,R} \).
Note that taking the $\bar{h}$ given by this corollary and periodizing it after reflection allows to show that $\inf_{\tilde{A}} \mathcal{W}_\eta$ has a minimizing sequence made of periodic vector fields.

**Proof of Proposition 5.10.** In view of Corollary 5.15, it suffices to bound from below $\lim \sup_{R \to +\infty} \int_{K_R} |\nabla \bar{h}_\eta|^2$. But the property of good separation of the points associated to $\bar{h}$ allows to apply Lemmas 5.6 and 3.6, which combined together yield that
\[
\int_{K_R} |\nabla \bar{h}_\eta|^2 - \#(\Lambda \cap K_R)(\kappa_d g(\eta) + \gamma_2 1_{d=2}) \geq -C \#(\Lambda \cap K_R),
\]
for some $C$ depending only on $d$. Since $\#(\Lambda \cap K_R) = |K_R|$, this entails
\[
\int_{K_R} |\nabla \bar{h}_\eta|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \geq -C.
\]
which implies the claimed lower bound. Of course, letting $\eta \to 0$ of course implies the same lower bound for $\mathcal{W}$, and this completes the proof.

Let us now give a sketch of the proof of Proposition 5.11, whose argument is based on [Lie2].

**Proof.** We treat the case of simple points, the case of multiple points can be ruled out in the same way as a limiting case of simple points.

Let $\Lambda$ be the set of points, and let us look a particular point of $\Lambda$, which we can assume, up to translation, to be the origin. For simplicity of the presentation, we will neglect the boundary effects and do as if the configuration lived in the whole space $\mathbb{R}^d$. We wish to show that if $\Lambda \setminus \{0\}$ contains a point $x$ very close to $0$, then the point $x$ can be moved away from $0$ to a point $y$ while decreasing the energy.

Let $U$ be the potential generated by all the points of $\Lambda$ except $x$ (in particular $U$ is regular in a neighborhood of $x$) and by the background distribution (which here is constant, but weaker assumptions would suffice). Suppose that we modify the configuration by moving $x$ to $y$, and let $\bar{h}$ be the perturbation induced on the electrostatic potential, i.e.
\[
\begin{cases}
-\Delta \bar{h} = c_d \left( \delta_y^{(n)} - \delta_x^{(n)} \right) & \text{in } K_R \\
\bar{h} = 0 & \text{on } \partial K_R.
\end{cases}
\]

We want to estimate the energy of the perturbed configuration. Several integrations by parts allow to see that
\[
\int_{K_R} |\nabla (h + \bar{h})|^2 - \int_{K_R} |\nabla h|^2 = \int_{K_R} |\nabla \bar{h}|^2 + 2c_d \int_{K_R} (\delta_y^{(n)} - \delta_x^{(n)}) \bar{h}
= 2c_d D_R(\delta_0^{(n)}, \delta_0^{(n)}) - 2c_d \int_{K_R} g_R * \delta_x^{(n)} \delta_y^{(n)} + 2c_d \int_{K_R} (\delta_y^{(n)} - \delta_x^{(n)}) (U + g_R * \delta_x^{(n)})
= 2c_d \int_{K_R} U(\delta_y^{(n)} - \delta_x^{(n)}),
\]
(5.59)
where \( D_R \) and \( g_R \) are respectively the Coulombic quadratic form (analogue to \( D \)) and the Coulomb kernel with zero boundary condition on \( \partial K_R \). Thus, if the configuration minimizes the energy, then both sides must always be nonnegative hence \( x \) must be at a local (even global) minimum of \( U \). This is true at least formally for \( \eta = 0 \), when we do not smear out the charges, but can be adapted to smeared out charges with errors that become negligible as \( \eta \) gets small. On the other hand, the potential \( U \) may be decomposed as

\[
U = U^{\text{ext}} + U^{\text{int}}
\]

where

\[
U^{\text{int}} = c_d \Delta^{-1} \left( \delta_0^{(\eta)} - 1_{B(0,2r_0)} \right)
\]

is the potential created by the singular charge at 0 and by the background distribution in the ball, and

\[
U^{\text{ext}} = c_d \Delta^{-1} \left( \sum_{p \in \Lambda \setminus \{x, 0\}} \delta_p^{(\eta)} - 1_{B(0,2r_0)^c} \right)
\]

is the potential generated by all the other charges, without \( x \) (and the background distribution). We may then observe that \( U^{\text{ext}} \) is super-harmonic in the ball \( B(0,2r_0) \) and thus achieves its minimum at some point \( \bar{x} \) that belongs to the boundary \( \partial B(0,2r_0) \). Moreover, \( U^{\text{int}} \) is radial and explicitly computable, and if \( r_0 \) is small enough one can check that \( U^{\text{int}}(r) \) is decreasing and thus \( U^{\text{int}} \) achieves its minimum over \( B(0,2r_0) \) on the boundary \( \partial B(0,2r_0) \). But then if \( x \in B(0,2r_0) \) and \( x \neq 0 \), \( x \) can be moved to \( y = \bar{x} \), this decreases \( U \), hence in view of (5.59), this decreases the energy (how much it can decreased can be better estimated, and this quantitative version of the argument allows to adapt the proof to the case of a bounded set and with \( \eta \) nonzero). This shows the desired result: if 0 and \( x \) (hence two arbitrary points in the configuration) are not separated by a distance \( 2r_0 \) depending only on \( d \), then the energy can be decreased. \( \square \)
6 Deriving $\mathcal{W}$ as the large $n$ limit

Our goal in this chapter is to pass to the limit $n \to \infty$ in the results obtained in Chapter 4, say starting from Proposition 4.5 and (4.20), in order to extract $\mathcal{W}$ as a limiting energy. The main task is to obtain a lower bound in the limit $n \to \infty$, which is expressed in terms of an average of $\mathcal{W}$ with respect to a suitable measure that encodes all the possible blow-up profiles. This is accomplished via a general method which can be formulated abstractly, and which we start by presenting.

6.1 Lower bound for 2-scales energies

In this section we present the abstract framework which serves to prove lower bounds on energies containing two scales (one much smaller than the other). The question is to deduce from a $\Gamma$-convergence (as defined in section 2.1) result at a certain scale a statement at a larger scale. The framework can thus be seen as a type of $\Gamma$-convergence result for 2-scale energies. The lower bound is expressed in terms of a probability measure, which can be seen as a Young measure on profiles (i.e. limits of the configuration functions viewed in the small scale). The method is similar in spirit to that of Alberti-Müller [AlMu], where they introduce what they call “Young measures on micropatterns,” but differs a bit, in particular in the fact that it is based on the use of Wiener’s multiparameter ergodic theorem, following a suggestion of S. R. S. Varadhan.

Let us first give a rough idea of the type of situation we wish to consider. Let us assume we want to bound from below an energy which is the average over large (as $\varepsilon \to 0$) domains $\Omega_\varepsilon$ of some nonnegative energy density $e_\varepsilon(u)$, defined on a space of functions $X$ (functions over $\mathbb{R}^n$), $\int_{\Omega_\varepsilon} e_\varepsilon(u(x)) \, dx$, and we know the $\Gamma$-liminf behavior of $e_\varepsilon(u)$ on small (i.e. here, bounded) scales — here the two scales are the finite scale 1 and the large scale corresponding to the diameter of the large domain $\Omega_\varepsilon$. By this we mean that we know how to obtain bounds from below independent of $\varepsilon$, say for example we can prove that

$$\liminf_{\varepsilon \to 0} \int_{B_R} e_\varepsilon(u) \, dx \geq \int_{B_R} e(u) \, dx. \quad (6.1)$$

However we cannot always directly apply such a knowledge to obtain a lower bound on the average over large domains: A natural idea is to cut the domain $\Omega_\varepsilon$ into boxes of fixed size $R$, to obtain lower bounds on each box (say of the type (6.1)) and add them together. By doing so, we may lose some information on the behavior of the function on the boundary of the boxes, which would be necessary to obtain a nontrivial lower bound. Moreover, we only get a lower bound by a number (related to the minimal value that the lower bound can take, say e.g.
min_u \int_{B_n} e(u) \, dx), while we would prefer instead a lower bound which is still a function depending on the u’s, i.e. on limits of the configuration u_\varepsilon. This is achieved by using the multiparameter ergodic theorem, as we shall now describe.

Let us turn to more precise statements. Let Ω be a compact set of positive measure in \(\mathbb{R}^d\), satisfying

\[
\lim_{\varepsilon \to 0} \frac{|(\Omega + \varepsilon x) \triangle \Omega|}{|\Omega|} = 0
\]

(6.2)

(where \(\triangle\) denotes the symmetric difference between sets). For each \(\varepsilon\), let \(f_\varepsilon(x,u)\) be a functional depending on \(x\), defined on a space of functions on \(\mathbb{R}^d\), assumed to be a Polish space, and denoted \(X\). We require \(f_\varepsilon\) to be measurable functions on \([\Omega \times X\).

Example 6.1. The function given by

\[
f_\varepsilon(x,u) = \int p(x)e_\varepsilon(u)\chi(y)dy
\]

where \(\chi\) is a cut-off function supported in \(B(0,1)\), \(e_\varepsilon\) is the energy density, and \(p\) is a function on \(\Omega\). The function \(p\) can be interpreted as a weight depending on \(x\), if \(p\) is constant then the functionals \(f_\varepsilon(x,\cdot)\) do not depend on \(x\).

We denote by \(\theta_\lambda\) the action of \(\mathbb{R}^d\) on the space \(X\) by translations, i.e. \(\theta_\lambda u = u(\lambda + \cdot)\) (it could be a more general action, but for the applications we have in mind, the action of translations is really what we need), and we require that \((\lambda,u) \mapsto \theta_\lambda u\) is continuous with respect to each variable. We also define the following groups of transformations on \(\mathbb{R}^d \times X\) :

\[
T_\varepsilon^\lambda(x,u) = (x + \varepsilon \lambda, \theta_\lambda u), \quad T_\lambda(x,u) = (x, \theta_\lambda u).
\]

We assume we are looking at a global energy of the form

\[
F_\varepsilon(u) = \int_{\Omega} f_\varepsilon(x,\theta_\varepsilon^\cdot u)dx.
\]

Example 6.2. If the local functional \(f_\varepsilon(x,u)\) is given by

\[
f_\varepsilon(x,u) = \int_{y \in \mathbb{R}^d} e_\varepsilon(u)\chi(y)dy
\]

where \(\chi(y)\) is a cut-off function of integral 1 supported in say \(B(0,1)\), and \(e_\varepsilon\) is the local energy density (this is the simpler case of Example 6.1 where \(f_\varepsilon(x,u)\) does
6.1 Lower bound for 2-scales energies

not depend on \(x\), then, with the previous definition, \(F_\varepsilon\) is equal to
\[
F_\varepsilon(u) = \int_{x \in \Omega} \left[ \int_{y \in \mathbb{R}^d} \chi(y)e_\varepsilon\left(u\left(\frac{x}{\varepsilon} + y\right)\right) dy \right] dx
\]
\[
= \varepsilon^d \int_{y \in \mathbb{R}^d} \int_{\frac{1}{\varepsilon} \Omega + y} \chi(y)e_\varepsilon(u(z)) dy dz
\]
\[
\approx \frac{\varepsilon^d}{|\Omega|} \left[ \int y \chi(y) dy \right] e_\varepsilon(u(z)) dz = \int_{\frac{1}{\varepsilon} \Omega} e_\varepsilon(u(z)) dz.
\]

The first equality is simply a change of variables \(z = \frac{x}{\varepsilon} + y\). Between the second and the third line, we note that the sets \(\frac{1}{\varepsilon} \Omega + y\) over which we integrate are almost constant: \(\frac{1}{\varepsilon} \Omega\) is of size \(1 \gg 1\) and we translate it by a small \(y \in B(0, 1)\). Therefore, an application of Fubini’s theorem and the use of (6.2) allow us to exchange the integration over \(y \in \mathbb{R}^d\) and the one over \(z \in \frac{1}{\varepsilon} \Omega\). Writing \(\int_{\frac{1}{\varepsilon} \Omega} e_\varepsilon(u(z)) dz\) in this fashion can be seen as a way to use a smooth partition of unity.

We will make the following assumptions:

(i) (bound from below) The functionals \(f_\varepsilon\) are bounded below by a constant independent of \(\varepsilon\) (for convenience we suppose, up to adding a constant, that \(f_\varepsilon \geq 0\)).

(ii) (coercivity and \(\Gamma\)-liminf) There exists a nonnegative measurable function \(f\) on \(\Omega \times X\), such that the following holds: if the quantities
\[
\int_{K_R} f_\varepsilon(T_\varepsilon^x(x_\varepsilon, u_\varepsilon)) d\lambda
\]
are bounded (when \(\varepsilon \to 0\)) for any \(R\), then \((x_\varepsilon, u_\varepsilon)\) has a convergent subsequence, converging to some \((x, u)\) and
\[
\liminf_{\varepsilon \to 0} f_\varepsilon(x_\varepsilon, u_\varepsilon) \geq f(x, u).
\]

The next step is to define what we announced as “Young measures on profiles.” For \(u\) in \(X\), we let \(P_\varepsilon\) be the probability measure on \(\Omega \times X\) obtained by pushing forward the normalized Lebesgue measure on \(\Omega\) by the map
\[
\Omega \to \Omega \times X
\]
\[
x \mapsto (x, \theta_{\frac{x}{\varepsilon}} u).
\]
It is equivalent to define \(P_\varepsilon\) as the probability measure such that for any \(\Phi \in C^0(\Omega \times X)\):
\[
\int \Phi(x, v) dP_\varepsilon(x, v) = \int_{\Omega} \Phi(x, \theta_{\frac{x}{\varepsilon}} u) dx.
\]  

(6.4)

We are thus considering the probability measures on the translates of the blow-ups of a given function \(u\) with the average obtained by centering the blow-up uniformly over the points of \(\Omega\). Formally one can write:
\[
P_\varepsilon = \int_{\Omega} \delta\left(x, \theta_{\frac{x}{\varepsilon}} u\right) dx.
\]
One can also be more precise by viewing \( u \mapsto P_\varepsilon \) as an embedding

\[
\left\{ \begin{array}{l}
\phi_\varepsilon : X \to \mathcal{P}(\Omega \times X) \\
u \mapsto \int_\Omega \delta_{(x, \theta_\varepsilon u)} \, dx,
\end{array} \right. \tag{6.5}
\]

where \( \mathcal{P}(S) \) denotes the space of Borel probability measures on \( S \). Note that the first variable \( x \) is just there to keep the memory of the blow-up center. The first marginal of \( P_\varepsilon \) is always equal to the normalized Lebesgue measure on \( \Omega \), regardless of the function \( u \). Also, the probability here is that of an analyst: the embedding \( \phi_\varepsilon \) is completely deterministic.

If \( P_\varepsilon = \phi_\varepsilon(u_\varepsilon) \) for a sequence of functions \( u_\varepsilon \) has a limit as \( \varepsilon \to 0 \), that limit can be seen as a Young measure, but encoding the whole blow-up profiles \( u = \lim_{\varepsilon \to 0} \theta_\varepsilon u_\varepsilon \) rather than only the limiting values of \( u_\varepsilon \) at \( x \), as is the case with the usual definition of Young measures (for which we refer to [Eva]). For example, if the functions \( u \) represent distributions of points, and if these form a lattice packed at scale \( \varepsilon \), the result is the average over a fundamental domain of the lattice “seen” from every possible origin. In a more general situation, \( P \) encodes the respective weights of the possible point patterns that emerge locally. One could imagine for example in dimension 2 a probability with weight \( p \) on triangular lattice configurations and weight \( 1 - p \) on square lattice configurations.

By definition of \( P_\varepsilon \), we can rewrite the global energy \( F_\varepsilon \) as the integral of the local energy \( f_\varepsilon \) with respect to \( P_\varepsilon \):

\[
F_\varepsilon(u_\varepsilon) = \int_\Omega f_\varepsilon(x, \theta_\varepsilon u_\varepsilon) \, dx = \int_{\Omega \times X} f_\varepsilon(x, v) dP_\varepsilon(x, v). \tag{6.6}
\]

Now, if we are able to find a limit \( P \) to the probability measures \( P_\varepsilon \) as \( \varepsilon \to 0 \), we may hope to write

\[
\liminf_{\varepsilon \to 0} \int f_\varepsilon \, dP_\varepsilon \geq \int f \, dP
\]

where \( f \) is given by assumption (ii). This will indeed hold and is reminiscent of Fatou’s lemma (indeed the sequence \( \{f_\varepsilon\} \) is, by assumption, bounded below). The last step is to combine this with the multiparameter ergodic theorem of Wiener (see [Bec]), whose statement we recall:

**Theorem 6.1** (Multiparameter ergodic theorem). Let \( X \) be a Polish (complete separable metric) space with a continuous \( d \)-parameter group \( \Theta_\lambda \) acting on it. Assume \( P \) is a \( \Theta \)-invariant probability measure on \( X \). Then for all \( f \in L^1(P) \), we have

\[
\int f(u) \, dP(u) = \int f^*(u) \, dP(u)
\]

where

\[
f^*(u) := \lim_{R \to +\infty} \frac{1}{\mu(K_R)} \int_{K_R} f(\theta_\lambda u) \, d\lambda \quad P\text{-a.e.}
\]

We may replace the cubic domains \( K_R \) by any family of reasonable shapes, such as balls, etc (more precisely a Vitali family, see [Bec] for the conditions).
Let us now give the statement of the abstract result. It originally appeared in [SS6] in the case where the energy density does not depend on the blow-up center \( x \), and was then generalized in [SS7].

**Theorem 6.2** (Lower bound for two-scale energies [SS7]). Assume \( \Omega, X, \{\theta_\lambda\}, \{F_\epsilon\}_\epsilon \) are as above and satisfy assumptions (i)–(ii). Assume \( \{u_\epsilon\}_\epsilon \), a family of elements of \( X \), is such that \( \{F_\epsilon(u_\epsilon)\}_\epsilon \) is bounded, and let \( P_\epsilon = \phi_\epsilon(u_\epsilon) \). Then

1. Up to extraction of a subsequence, \( \{P_\epsilon\}_\epsilon \) converges weakly in the sense of probabilities, to some probability measure \( P \in \mathcal{P}(\Omega \times X) \), whose first marginal is the normalized Lebesgue measure on \( \Omega \).
2. The limit \( P \) is \( T_\lambda \)-invariant.
3. \( P \)-almost every point \((x,u)\) is a limit of \((x_\epsilon, \theta_\lambda u_\epsilon)\). (Thus \( P \) is indeed an average over possible local limits.)
4. The following \( \liminf \) holds:

\[
\liminf_{\epsilon \to 0} F_\epsilon(u_\epsilon) \geq \int f(x,u) dP(x,u) = \int f^*(x,u) dP(x,u)
\]

with

\[
f^*(x,u) := \lim_{R \to +\infty} \int_{K_R} f(T_\lambda(x,u)) d\lambda = \lim_{R \to +\infty} \int_{K_R} f(x,\theta_\lambda u) d\lambda.
\]

We now indicate the ingredients of the proof (details can be found in [SS6,SS7]).

**Proof.** 1. The main point is to show that \( \{P_\epsilon\}_\epsilon \) is tight, i.e. for any \( \eta > 0 \) there exists a compact set \( K_\eta \) such that \( P_\epsilon(K_\eta) \geq 1 - \eta \) for small \( \epsilon \). This comes as a consequence of the assumption that \( \{F_\epsilon(u_\epsilon)\}_\epsilon \) is bounded and the coercivity assumption (ii) on the functionals. The fact that the first marginal is the normalized Lebesgue measure is obvious since it is true for each \( P_\epsilon \) and thus remains true in the limit.

2. The invariance by \( T_\lambda \) is a straightforward consequence of the definition of \( P_\epsilon \). Consider a test-function \( \Phi \in C^0(\Omega \times X) \) and \( \lambda \in \mathbb{R}^d \). On the one hand:

\[
\lim_{\epsilon \to 0} \int_{\Omega} \Phi(x,\theta_\lambda x u_\epsilon) dx = \lim_{\epsilon \to 0} \int_{\Omega} \Phi(x,\theta_\lambda u_\epsilon) dx
\]

because \( \frac{x}{\epsilon} + \lambda \approx \frac{x}{\epsilon} \) for any \( \lambda \) fixed when \( \epsilon \) goes to zero (this uses the assumption (6.2)). But on the other hand, for any \( \lambda \), we have by definition of \( P_\epsilon \),

\[
\lim_{\epsilon \to 0} \int_{\Omega} \Phi(x,\theta_\lambda x u_\epsilon) dx = \lim_{\epsilon \to 0} \int_{\Omega} \Phi(x,\theta_\lambda u_\epsilon) dP_\epsilon(x,u) = \int_{\Omega} \Phi(x,\theta_\lambda u_\epsilon) dP(x,u).
\]

We deduce that we must have, for all continuous \( \Phi \),

\[
\int_{\Omega} \Phi(x,\theta_\lambda u) dP = \int_{\Omega} \Phi(x,u) dP,
\]

(6.8)
Deriving $\mathcal{W}$ as the large $n$ limit

which exactly means that $P$ is $T_\lambda$-invariant.

3. This is a rather direct consequence of the definition of $P_\varepsilon$.

4. This is a result that uses the fine topological information provided by assumptions (ii), and combines it with the weak convergence of $P_\varepsilon$ to $P$, assumption (i) and Fatou’s lemma, cf. [SS6, Lemma 2.2].

As desired, this result provides a lower bound on functionals of the type $F_\varepsilon(u_\varepsilon)$, which is expressed in terms of the probability $P$, i.e. in terms of the limits of $u_\varepsilon$. As a corollary, it implies the weaker result of lower bound of $F_\varepsilon$ by a number:

$$\lim_{\varepsilon \to 0} \inf F_\varepsilon(u_\varepsilon) \geq \int_\Omega \inf_u f^*(x,u)dP(x,u) = \int_\Omega \inf_u f^*(x,u)dx.$$ (6.9)

The minimization of the function $f^*$ is similar to a “cell problem” in homogenization (cf. e.g. [BraDe2]).

Once this result is proved, it remains to show, if possible, that such a lower bound is sharp, which requires constructing a family $\{u_\varepsilon\}$ such that

$$\limsup_{\varepsilon \to 0} F_\varepsilon(u_\varepsilon) \leq \int_\Omega \inf f^*(x,\cdot))dx.$$ (6.10)

This certainly requires at least that the $\Gamma$-liminf relation in assumption (i) be also a $\Gamma$-limsup, i.e. that there exist recovery sequences. This really depends on the specifics of the local functionals. If (6.10) can be shown, then, just as in Proposition 2.6 comparing (6.9) and (6.10) implies that if $u_\varepsilon$ minimize $F_\varepsilon$ for every $\varepsilon$ and $\min F_\varepsilon$ is bounded, then letting $P$ be as in Theorem 6.2 we must have

$$P - a.e. (x,u), \ u \text{ minimizes the local functional } f^*(x,\cdot).$$

We will next see how to apply this abstract result in the context of the Coulomb gas Hamiltonian. It has also been used for vortices in Ginzburg-Landau in [SS6], as we will see in Chapter 10 and droplets in the Ohta-Kawasaki model [GMS2]. In all these cases, we were able to conclude because the corresponding upper bound (6.10) turned out to be provable.

6.2 Next order lower bound for the Coulomb gas Hamiltonian

6.2.1 Assumptions

To get to the final results of this chapter, we will make the following assumptions on $V$ :

1. the strongest ones made in Chapter 2 i.e. that $V$ is continuous and satisfies (A3)–(A4). This in particular guarantees that the equilibrium measure $\mu_0$ exists and has compact support.
2. \textbf{(A5)} The support $\Sigma$ of the equilibrium measure has a $C^1$ boundary.

3. \textbf{(A6)} The equilibrium measure $\mu_0$ has an $L^\infty$ density which is bounded below of class $C^1$ on its support:

$$
\mu_0(x) = m_0(x)1_\Sigma(x)dx
$$

(6.11)

where $\Sigma$ is the support of $\mu_0$ and $m_0 \in C^1(\Sigma) \cap L^\infty(\mathbb{R}^d)$ is its density, which satisfies

$$
0 < m_0 \leq m_0 \leq \overline{m}.
$$

(6.12)

Again, by abuse of notation, we will confuse $m_0(x)$ and $\mu_0(x)$.

If $V$ is smooth enough, these assumptions are sufficiently generic. They are nonempty: an easy example is the case when $V$ is a multiple of $|x|^2$ and $\mu_0$ is a multiple of the characteristic function of a ball (see Example 2 in Chapter 2) — in fact any $V$ positive quadratic works as well. Recall also that from Proposition 2.22 when $V$ is $C^2$ $\mu_0$ is a measure with density $m_0 = (\frac{1}{2}\Delta V 1_\omega) \in L^\infty$, thus if $\Delta V$ is bounded below by a positive constant, (6.12) is satisfied. If in addition $V$ is $C^3$ on $\Sigma$, then $\mu_0$ is $C^1$ in $\Sigma$ and (A6) is fully satisfied. This strong assumption is assumed mostly for convenience, to simplify our upper bound construction. For the lower bound, the assumption that $\mu_0 \in C^0(\Sigma)$ (and probably even less) suffices. Note that when (A6) holds, by continuity of $\Delta V$, $\Sigma$ and the coincidence set $\omega$ must coincide.

The assumption (A5) can be investigated in light of the regularity theory for the obstacle problem, for which $C^1$ regularity of the boundary of the coincidence set is generic in some sense in dimension 2 [Sc, Mon], or is true if the coincidence set is convex. Note also that a result [KN, Isa] shows that if the boundary of the coincidence set is $C^1$, it is in fact analytic. Again, weaker conditions should suffice.

6.2.2 Lower bound

As already mentioned, we now return to Proposition 4.5 in order to extract $\mathcal{W}$ as a limiting lower bound, using the abstract framework of Section 6.1.

In view of the results of Proposition 3.4, in order to bound from below $H_n$ at the next order, it suffices to bound from below $\mathcal{W}(\nabla h'_n, 1_\Sigma)$ where $h'_n$ is as in (3.24). Chapter 4 provides us with the first lower bound (4.20), which is our starting point. It then suffices to bound from below $\frac{1}{\eta} \int_{\mathbb{R}^d} |\nabla h'_{n, \eta}|^2$, where $h'_{n, \eta}$ is given by (4.17). This will be done according to the scheme of Section 6.1. We first consider $\eta$ as fixed and let $n \to \infty$, and later let $\eta \to 0$. The setup to use the abstract framework is to take $\Omega = \Sigma$, $X = L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d)$ for some $q < \frac{d}{d-1}$, and $\varepsilon = n^{-1/d}$. Assumption (A5) ensures in particular that the condition (6.2) is satisfied.

The main lower bound result that we obtain this way is

**Theorem 6.3** (Lower bound at next order for the Coulomb gas Hamiltonian). Assume the equilibrium measure $\mu_0$ exists and satisfies (A5)–(A6). For any
$x_1, \ldots, x_n \in \mathbb{R}^d$, let $h_n' = g \ast \left( \sum_{i=1}^n \delta_{x_i'} - \mu_0' \right)$ as in \eqref{3.24}. Let $P_n \in \mathcal{P}(\Sigma \times X)$ be the push-forward of the normalized Lebesgue measure on $\Sigma$ by

$$x \mapsto (x, \nabla h_n'(n^{1/d}x + \cdot)).$$

Then, up to extraction of a subsequence, $P_n$ converges weakly in the sense of probabilities to a probability measure $P \in \mathcal{P}(\Sigma \times X)$ such that

(i) $P$ is translation-invariant, and its first marginal is the normalized Lebesgue measure on $\Sigma$.

(ii) For $P$-almost every $(x, E)$, $E$ belongs to the class $\mathcal{A}_{\mu_0(x)}$.

(iii) We have the following $\lim \inf$ inequality:

$$\lim \inf_{n \to +\infty} n^{2/d-2} \left( H_n(x_1, \ldots, x_n) - n^2 I(\mu_0) - 2n \sum_{i=1}^n \zeta(x_i) + \left( \frac{n}{2} \log n \right) 1_{d=2} \right) \geq \tilde{W}(P), \tag{6.13}$$

where $\tilde{W}$ is defined over the set of probability measures $P \in \mathcal{P}(\Sigma \times X)$ satisfying (i) and (ii) by

$$\tilde{W}(P) := \frac{|\Sigma|}{cd} \int \mathcal{W}(E) \, dP(x, E). \tag{6.14}$$

This result was proven in this form in \cite{RouSe}. The same also holds with $\mathcal{W}$ replaced by $W$ in dimension $d = 2$, as was previously proven in \cite{SS7}, and also in dimension $d = 1$ in \cite{SS8}, with $\frac{n}{2} \log n$ replaced by $n \log n$, as in Proposition 3.5.

One may guess the value of the minimum of $\tilde{W}$ on its domain of definition: by property (i) on $P$, we have

$$\min \tilde{W} \geq \frac{|\Sigma|}{cd} \int \left( \min_{\mathcal{A}_{\mu_0(x)}} \mathcal{W} \right) dP(x, E) = \frac{1}{cd} \int_{\Sigma} \left( \min_{\mathcal{A}_{\mu_0(x)}} \mathcal{W} \right) dx, \tag{6.15}$$

and by the scaling relations \eqref{5.14} and \eqref{5.15}, we thus get

$$\min \tilde{W} \geq \xi_d := \frac{1}{cd} \int_{\Sigma} \min_{\mathcal{A}_{\mu_0(x)}} \mathcal{W} \, dx$$

$$= \begin{cases} \frac{1}{cd} \left( \int_{\Sigma} \mu_0^{-2/d}(x) \, dx \right) \min_{\mathcal{A}_1} \mathcal{W} & \text{for } d \geq 3 \\ \frac{1}{2\pi} \min_{\mathcal{A}_1} \mathcal{W} - \frac{1}{2} \int_{\Sigma} \mu_0(x) \log \mu_0(x) \, dx & \text{for } d = 2. \end{cases} \tag{6.16}$$

It turns out, as we will see below, that these inequalities are equalities. Dropping the term $\sum \zeta(x_i)$ which is always nonnegative, Theorem 6.3 has the following
Corollary 6.3. We have
\[
\liminf_{n \to +\infty} n^{2/d-2} \left( \min H_n - n^2 I(\mu_0) + \left( \frac{n}{2} \log n \right) \mathbf{1}_{d=2} \right) \geq \xi_d.
\]

Proof of the theorem. As announced, we apply the abstract framework of Section 6.1 for fixed \( \eta \). We will need the following notation: given a gradient vector field \( E = \nabla h \) satisfying a relation of the form
\[
- \text{div} E = c_d \left( \sum_{\mu \in \Lambda} \sum_{p \in A} N_p \delta_p - \mu(x) \right) \text{ in } \mathbb{R}^d,
\]
whether an element of \( \bar{A}_m \), or the gradient of a potential defined by (4.17), we define \( E_{\eta} \) to be as in (5.5), and we denote \( \Phi_{\eta} \) for the map \( E \mapsto E_{\eta} \), which to a vector field corresponding to singular charges assigns the vector field corresponding to smeared out charges.

Let us define \( P_{n,\eta} \) as the push-forward of the normalized Lebesgue measure on \( \Sigma \) by the map
\[
x \mapsto (x, \nabla h'_{\eta}(n^{1/d}x + \cdot)).
\]
In other terms, \( P_{n,\eta} \) is the push-forward of \( P_n \) by \( \Phi_{\eta} \). Then, we take \( \chi \) to be a smooth nonnegative cut-off function supported in \( B(0,1) \) and of integral 1, and set
\[
f_n(x,E) = \begin{cases} 
\int \chi(y)|E|^2(y)dy & \text{if } E = \nabla h'_{\eta}(n^{1/d}x + \cdot) \\
+\infty & \text{otherwise.}
\end{cases}
\]
This gives the “local” energy at the small scale. The definition ensures that we only consider a class of vector fields that are of the interesting form.

We then let \( F_n(E) \) be given, as in Theorem 6.2, by
\[
F_n(E) = \int_E f_n(x,\theta_{n^{1/d}}E)dx
\]
and we may observe that by Fubini’s theorem and a change of variables,
\[
F_n(E) = \frac{1}{|\Sigma|} \int_{\mathbb{R}^d} \int_{\Sigma} \chi(y)|\nabla h'_{n,\eta}(n^{1/d}x + y)|^2 dxdy \\
\leq \frac{1}{\pi |\Sigma|} \int_{\mathbb{R}^d} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}(z)|^2 \chi(y) dydz.
\]
Since \( \int \chi = 1 \), it follows that
\[
\frac{1}{n} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 \geq |\Sigma| F_n(E)
\]
so in order to bound from below \( \frac{1}{n} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 \) as desired, it does suffice to bound from below \( F_n \), which will be done via Theorem 6.2.

To do so, we have to check the three assumptions (i)–(ii) stated in Section 6.1.
First, it is true that the energies \( f_n \) are nonnegative, this gives the condition (i). To check condition (ii), we use the following lemma.
Lemma 6.4 (Weak compactness of local electric fields).
Let \( h'_{n,q} \) be as above, and let \( \nu'_n = \sum_{i=1}^n \delta_{x'_i} \). Assume that for every \( R > 1 \) and for some \( \eta \in (0, 1) \), we have

\[
\sup_n \int_{K_R} |\nabla h'_{n,q}(n^{1/d}x_n + \cdot)|^2 \leq C_{\eta,R}, \tag{6.19}
\]

and \( x_n \to x \in \mathbb{R}^d \) as \( n \to \infty \) (a sequence of blow-up centers). Then \( \{\nu'_n(n^{1/d}x_n + \cdot)\}_n \) is locally bounded and up to extraction converges weakly as \( n \to \infty \) in the sense of measures to

\[
\nu = \sum_{p \in \Lambda} N_p \delta_p
\]

where \( \Lambda \) is a discrete set of \( \mathbb{R}^d \) and \( N_p \in \mathbb{N}^+ \). In addition, there exists \( E \in L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d) \), \( q < \frac{d}{d-1} \), \( E_\eta \in L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d) \), with \( E_\eta = \Phi_\eta(E) \), such that, up to extraction, as \( n \to \infty \),

\[
\nabla h'_{n,q}(n^{1/d}x_n + \cdot) \to E \text{ weakly in } L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d) \text{ for } q < \frac{d}{d-1}
\]

\[
\nabla h'_{n,q}(n^{1/d}x_n + \cdot) \to E_\eta \text{ weakly in } L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d). \tag{6.20}
\]

Moreover \( E \) is the gradient of a function \( h \), and if \( x \notin \partial \Sigma \), we have

\[
-\Delta h = c_d(\nu - \mu_0(x)) \quad \text{in } \mathbb{R}^d \tag{6.22}
\]

hence \( E \in \mathcal{A}_{\mu_0(x)} \).

Proof. First, from (6.19) and (3.24), exactly as in the proof of (5.19) we have that for some \( t \in (R-1, R) \), for every \( n \),

\[
\left| \int_{K_t} \sum_i \delta_{x'_i} (n^{1/d}x_n + x') - \int_{K_t} \mu_0(x_n + n^{-1/d}x') \, dx' \right| \leq C_{\eta,R},
\]

for some constant depending only on \( \eta \) and \( R \). It follows that, letting \( \nu'_n := \nu'_n(n^{1/d}x_n + \cdot) \), we have

\[
\nu'_n(K_{R-1}) \leq C_d\|\mu_0\|_{L^\infty} R^d + C_{\eta,R}.
\]

This establishes that \( \{\nu'_n\} \) is locally bounded independently of \( n \). In view of the form of \( \nu'_n \), its limit can only be of the form \( \nu = \sum_{p \in \Lambda} N_p \delta_p \), where \( N_p \) are positive integers and \( \Lambda \) is a discrete set in \( \mathbb{R}^d \).

Up to a further extraction we also have (6.21) by (6.19) and weak compactness of \( \nabla h'_{n,q} \) in \( L^2_{\text{loc}}. \) The compactness and convergence (6.20) follow from Lemma 4.9. The weak local convergences of both \( \nu'_n \) and \( \nabla h'_{n,q}(n^{1/d}x_n + \cdot) \) together with the continuity of \( \mu_0 \) away from \( \partial \Sigma \) (cf. (A0)), imply after passing to the limit in

\[
-\Delta h'_n(n^{1/d}x_n + \cdot) = c_d \left( \nu'_n - \mu_0(x_n + n^{-1/d}x') \right)
\]
(which is obtained by centering \(3.18\) around \(x_n\)) that \(E\) must be a gradient and that \(6.22\) holds. Finally \(E_0 = \Phi_\eta(E)\) because one may check that \(\Phi_\eta\) commutes with the weak convergence in \(L^q_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d)\) for the \(\nabla h_n^\prime(n^{1/d}x_n + \cdot)\) described above.

To check condition (ii), let us thus assume that

\[
\forall R > 0, \limsup_{n \to +\infty} \int_{K_R} f_n(T^n(x_n, Y_n)) \, d\lambda < \infty, \quad x_n \in \Sigma.
\]

By definition of \(f_n\), this condition is equivalent to

\[
\forall R > 0, \forall n \geq n_0, Y_n = \nabla h_n^\prime(\cdot) \quad \text{and} \quad \limsup_{n \to +\infty} \int \chi^* 1_{K_R} |Y_n|^2 < +\infty, \quad x_n \in \Sigma.
\]

This implies the assumption of Lemma 6.4. We may also assume, up to extraction of a subsequence that \(x_n \to x \in \Sigma\) (since \(\Sigma\) is compact). Applying Lemma 6.4, we have \(Y_n \rightharpoonup Y\) weakly in \(L^2_{\text{loc}}(\mathbb{R}^d, \mathbb{R}^d)\), and all the other results of the lemma. This weak convergence implies in particular that

\[
\liminf_{n \to +\infty} f_n(x_n, Y_n) \geq f(x, Y)
\]

\[
:= \begin{cases} 
\int \chi(y)|Y|^2(y) \, dy & \text{if } x \in \Sigma \setminus \partial \Sigma \text{ and } Y = \Phi_\eta(E) \text{ for some } E \in \bar{A}_{\mu(x)} \\
0 & \text{if } x \in \partial \Sigma \\
+\infty & \text{otherwise.}
\end{cases}
\]

This completes the proof that condition (ii) holds.

Theorem 6.2 then yields the convergence (up to extraction) of \(P_{n, \eta}\) to some \(P_\eta\), and, in view of \(6.18\),

\[
\liminf_{n \to +\infty} \int_{\mathbb{R}^d} |\nabla h_n^\prime|^2 \geq \liminf_{n \to +\infty} |\Sigma| F_n(E) \geq |\Sigma| \int f^*(x, Y) dP_\eta(x, Y) \quad (6.23)
\]

where

\[
f^*(x, Y) = \lim_{R \to +\infty} \int_{K_R} f(x, \theta_y Y) \, d\lambda.
\]

By definition of \(f\) and since \(\partial \Sigma\) is of Lebesgue measure 0 by (A5), for \(P_\eta\)-a.e. \((x, Y)\), we must have \(Y = \Phi_\eta(E)\) with some \(E \in \bar{A}_{\mu(x)}\). Pushing forward by \(\Phi_\eta^{-1}\), we get the convergence of \(P_n\) to \(P\) satisfying the first two stated properties. Moreover, applying Fubini’s theorem, we may write,

\[
f^*(x, Y) = \lim_{R \to +\infty} \int_{K_R} \chi * 1_{K_R} |Y|^2 \geq \lim_{R \to +\infty} \int_{K_R} |Y|^2
\]

where we used that \(\chi * 1_{K_R} \geq 1_{K_{R-1}}\). By definition of the push-forward, it follows from \(6.23\) that

\[
\liminf_{n \to +\infty} \int_{\mathbb{R}^d} |\nabla h_n^\prime|^2 \geq |\Sigma| \int \left( \lim_{R \to +\infty} \int_{K_R} |\Phi_\eta(E)|^2 \right) dP(x, E).
\]
Using that \( \int \mu_0 = 1 \), the fact that the first marginal of \( P \) is the normalized Lebesgue measure, that \( P \)-a.e., \( E \in \mathcal{A}_{\mu_0(x)} \), and the definition of \( \mathcal{W}_\eta \) (Definition 5.2), we deduce that

\[
\lim_{n \to +\infty} \frac{1}{n} \int_{\mathbb{R}^d} |\nabla h'_{n,\eta}|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \geq |\Sigma| \int \left( \lim_{R \to +\infty} \int_{K_R} |\Phi_\eta(E)|^2 - (\kappa_d g(\eta) + \gamma_2 1_{d=2}) \mu_0(x) \right) dP(x, E)
\]

\[
= |\Sigma| \int \mathcal{W}_\eta(E) dP(x, E).
\]

Inserting into (4.20), we find that

\[
\lim_{n \to \infty} n^{2/d-2} W(\nabla h'_n, 1_{\mathbb{R}^d}) \geq |\Sigma| \int \mathcal{W}_\eta(E) dP(x, E) - C\eta^2.
\]

Since \( \mathcal{W}_\eta \) is bounded below as seen in Proposition 5.10, we may apply Fatou’s lemma and take the \( \lim \inf \eta \to 0 \) on both sides, and we obtain, by definition of \( \mathcal{W} \), that

\[
\lim_{n \to \infty} n^{2/d-2} W(\nabla h'_n, 1_{\mathbb{R}^d}) \geq |\Sigma| \int \mathcal{W}(E) dP(x, E).
\]

Combining this with the result of Proposition 3.4, the conclusion (6.13) follows.

\[ \square \]

**Remark 6.5.** Note that for this lower bound, we do not really need the full strength of (A5)–(A6): it suffices that \( \mu_0 \) be continuous on its support, that \( \partial \Sigma \) has zero measure, and that \( \Sigma \) satisfies (6.2).

### 6.3 Upper bound and consequences for ground states

As already mentioned, to be useful, this lower bounds needs to be complemented by a corresponding upper bound, proving that the lower bound was sharp. As in \( \Gamma \)-convergence, this is accomplished by an explicit construction.

The following proposition states the result we can obtain. It shows that we can find some configuration of points for which the lower bound of Corollary 6.3 is sharp. Because we have in view the statistical mechanics problem as well, it will be useful to show that this is true not only for that configuration, but for a “thick enough” neighborhood of it.

**Proposition 6.6** (Upper bound at next order). For any \( \varepsilon > 0 \) there exists \( r_1 > 0 \) and for any \( n \) a set \( A_n \subset (\mathbb{R}^d)^n \) such that

\[
|A_n| \geq n! \left( \pi(r_1)^d / n \right)^n
\]

(6.24)
and for any \((y_1, \ldots, y_n) \in A_n\) we have
\[
\limsup_{n \to \infty} n^{2/d-2} \left( H_n(y_1, \ldots, y_n) - n^2 I(\mu_0) + \left( \frac{n}{2} \log n \right) 1_{d=2} \right) \leq \xi_d + \varepsilon, \tag{6.25}
\]
where \(\xi_d\) is defined by (6.16).

**Proof.** We sketch the main steps of the construction, which relies on pasting together vector fields obtained via the screening construction (cf. Proposition 5.12), more precisely those given by Corollary 5.15. That corollary was stated for hypercubes but it applies to hyperrectangles as well.

**Step 1.** We fix some large \(R\) and partition \(\Sigma'\) (the blown-up of the set \(\Sigma\)) into hyperrectangles \(K_i\) of sidelengths in \([R, 2R]\) and such that \(\int_{K_i} \mu_0' \in \mathbb{N}\). This is not very difficult to do, cf. [SS7, Lemma 7.5], it leaves however a (layer) region \(\Sigma'_{\text{bound}}\) near the boundary of \(\Sigma'\) which cannot exactly be partitioned into hyperrectangles. We let \(m_i = -\int_{K_i} \mu_0'\).

**Step 2.** We paste in each \(K_i\), copy of the \(\bar{h}\) given by Corollary 5.15, translated and rescaled by a factor \(m_i^{1/d}\), so that we have a solution to
\[
\begin{cases}
-\Delta h_i = c_d(\sum_p \delta_p - m_i) & \text{in } K_i \\
\frac{\partial h_i}{\partial \nu} = 0 & \text{on } \partial K_i,
\end{cases}
\]
and
\[
\int_{K_i} |\nabla (h_i)\eta|^2 - m_i(\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq \min_{A_{m_i}} W + o_R(1). \tag{6.26}
\]
Note that the rescaling factor does not degenerate by assumption (A6), that the total number of points in \(K_i\) is \(m_i|K_i| = \int_{K_i} \mu_0'\), and that the points are separated by a distance depending only on \(m\) in Assumption (A6), as provided by Corollary 5.15.

**Step 3.** Since \(m_i\) is not the desired weight \(\mu_0'\), we correct \(h_i\) by adding a solution to
\[
\begin{cases}
-\Delta u_i = c_d(m_i - \mu_0') & \text{in } K_i \\
\frac{\partial u_i}{\partial \nu} = 0 & \text{on } \partial K_i.
\end{cases}
\]
Thanks to assumption (A6), we know that \(\mu_0'\) varies slowly (more precisely \(|\nabla \mu_0'| \leq C n^{-1/d}\) at the scale considered), so \(\|m_i - \mu_0'\|_{L^\infty(K_i)}\) is small, which allows to prove by elliptic regularity estimates that \(u_i\) is small in a strong sense. We note that this is the point where we use the \(C^1\) regularity of \(\mu_0\) on its support, but that it could easily be replaced by a weaker statement showing slow variation (such as a Hölder continuity assumption). We then let \(E_i = \nabla h_i + \nabla u_i\) in each \(K_i\).

**Step 4.** We complete the construction by defining a vector field of the right form in the region \(\Sigma'_{\text{bound}}\) near the boundary of \(\Sigma'\). Because that region has a negligible volume, it is not important to use an approximate minimizer of \(W\), it
Deriving $\mathcal{W}$ as the large $n$ limit suffices to construct some vector field $E_{\text{bound}}$ associated to well-separated points, and satisfying

\[
\begin{cases}
-\text{div} E_{\text{bound}} = c_d(\sum_{p \in \Lambda_{\text{bound}}} \delta_p - \mu'_0) & \text{in } \Sigma'_{\text{bound}} \\
E_{\text{bound}} \cdot \nu = 0 & \text{on } \partial \Sigma'_{\text{bound}}.
\end{cases}
\]

and

\[
\int_{\Sigma'_{\text{bound}}} |E_{\text{bound}, \eta}|^2 - \# \Lambda_{\text{bound}}(\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq o(n),
\]

(6.27)

**Step 5.** We paste together the vector fields $E_i$ and $E_{\text{bound}}$ defined in all the regions that make up $\Sigma'$, extend them by 0 outside $\Sigma'$, and call the result $E$. Because the normal components of these vector fields are continuous across the interfaces between the regions, $E$ globally satisfies a relation of the form

\[-\text{div} E = c_d(\sum_{p \in \Lambda_n} \delta_p - \mu'_0) \text{ in } \mathbb{R}^d,\]

for a collection of points $\Lambda_n$ in $\Sigma'$ which are simple and well separated, and for which we can check that $\# \Lambda_n = n$ (since $\int_{\Sigma'} \mu'_0 = n \int \mu_0 = n$). The vector field $E$ is no longer a gradient, however we can keep the points of $\Lambda_n$ and define $h'_{n, \eta}$ associated to them via (3.24). Computing exactly as in Step 4 of the proof of Proposition 5.12. shows that

\[
\hat{R}^d |\nabla h'_{n, \eta}|^2 \leq \hat{R}^d |E_\eta|^2 \leq \sum_i |K_i| \min_{A_{m_i}} \mathcal{W} + \text{negligible terms},
\]

(6.28)

**Step 6.** We estimate the energy of the constructed configuration. Adding the contributions given by (6.26) and (6.27), and inserting into (6.28), we obtain that for $R$ and $\eta$ given, we have a configuration of points $\{x'_1, \ldots, x'_n\} = \Lambda_n$ (and a corresponding blown-down configuration $(x_1, \ldots, x_n)$) for which

\[
\int_{\mathbb{R}^d} |\nabla h'_{n, \eta}|^2 - n(\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq \sum_i |K_i| \min_{A_{m_i}} \mathcal{W} + \text{negligible terms}.
\]

Using a Riemann sum argument and the continuity of $m \mapsto \min_{A_m} \mathcal{W}$ which follows from (5.14)–(5.15), we conclude that

\[
\int_{\mathbb{R}^d} |\nabla h'_{n, \eta}|^2 - n(\kappa_d g(\eta) + \gamma_2 1_{d=2}) \leq \int_{\Sigma'} \left( \min_{A_{m_i}(\eta)} \mathcal{W} \right) dx + o_R(n).
\]

Since the points are well-separated, as soon as $\eta$ is small enough, we are in the case of equality of Proposition 4.5, which means that we may write (using also
that \( \zeta(x_i) = 0 \) by Definition 2.18 since all the \( x_i \)’s are in \( \Sigma \) )

\[
H_n(x_1, \ldots, x_n) \leq n^2 I(\mu_0) - \left( \frac{n}{2} \log n \right) 1_{d=2} \\
\quad + \frac{n^{2-2/d}}{c_d} \left( \frac{1}{n} \int_{\Sigma'} \min_{A_{\mu_0}(x)} W_\eta \, dx + o_R(1) \right).
\]

Taking \( R \) large enough and \( \eta \) small enough, and using the fact that \( \min_{A_{\mu_0}} W_\eta \to \min_{A_{\mu_0}} W \) as \( \eta \to 0 \), and by definition of \( \xi_d \), we can find \((x_1, \ldots, x_n)\) so that (6.25) holds, i.e. so that we have the desired right-hand side up to an error \( \epsilon \).

**Step 7.** The statement about the volume of the set \( A_n \) follows by noting that if \( y_i \in B(x_i, r_1 n^{-1/d}) \) for \( r_1 \) small enough (depending on \( r_0 \)) then \( y_1, \ldots, y_n \) are also well separated and we may perform the same analysis for \( H_n(y_1, \ldots, y_n) \), except with an additional error depending only on \( r_1 \) and going to 0 when \( r_1 \to 0 \), so which can be made \( < \epsilon \) by taking \( r_1 \) small enough. It is in addition clear that the set of such \( y_i \)’s has volume \( n!(r_1/n)^n \) in configuration space: the \( r_1/n \) term is the volume of the ball \( B(x_i, r_1 n^{-1/d}) \), it is raised to the power \( n \) because there are \( n \) points in the configuration, and multiplied by \( n! \) because permuting \( y_1, \ldots, y_n \) does not change the energy.

**Remark 6.7.** If we view these results as a \( \Gamma \)-convergence at next order of \( H_n \), then Theorem 6.3 provides the \( \Gamma \)-liminf relation, but Proposition 6.6 only provides the \( \Gamma \)-limsup relation at the level of minimizers, which is enough to conclude about these, but does not provide a full \( \Gamma \)-convergence result. To get one, we would need to construct recovery sequences for all probabilities \( P \) satisfying the first two properties of Theorem 6.3. This is technically more complicated, because such \( P \)’s need to be approximated by a single sequence, and it was accomplished in [SS7] for \( W \) in dimension 2. In the setting of \( \mathcal{W} \), we are limited by the fact that our screening procedure is written only for configurations with well-separated points.

Comparing with the lower bound result of Theorem 6.3, we immediately obtain the following result on the minimum and minimizers of the Coulomb gas Hamiltonian.

**Theorem 6.4 (Ground state energy expansion and microscopic behavior of minimizers).**

Assume \( V \) is such that the equilibrium measure exists and satisfies (A5)–(A6).

As \( n \to \infty \) we have

\[
\min H_n = n^2 I(\mu_0) - \left( \frac{n}{2} \log n \right) 1_{d=2} + n^{2-2/d} \xi_d + o(n^{2-2/d}),
\]

where \( \xi_d \) is as in (6.16), and it holds that \( \xi_d = \min \tilde{\mathcal{W}} \). In addition, if \((x_1, \ldots, x_n) \in (\mathbb{R}^d)^n \) minimize \( H_n \), letting \( h_n' \) be associated via (3.18) and \( P_n \) and \( P \) be as in Theorem 6.3, then \( P \) minimizes \( \tilde{\mathcal{W}} \) and for \( P \text{-a.e.} (x,E) \), \( E \) minimizes \( \mathcal{W} \) over \( \tilde{A}_{\mu_0}(x) \).
This result was obtained in RouSe for \( d \geq 2 \) and previously for \( d = 2 \) in SS7 with \( W \) instead of \( W \), which gives another proof that in that case \( \min_{A_1} W = \min_{A_1} \mathcal{W} \) (we have seen other ways of justifying this in Corollary 5.8). In SS8 the corresponding result for \( d = 1 \) is obtained:

\[
\min H_n = n^2 I(\mu_0) - n \log n + n \left( \frac{1}{2\pi} \min_{A_1} W - \int_\Sigma \mu_0(x) \log \mu_0(x) \, dx \right) + o(n) \tag{6.30}
\]

and we recall that in that case the value of the minimum of \( W \) is known, cf. Theorem 6.1.

These expansions of the minimal energy are to be compared to Theorem 2.2; we have obtained as announced a next order expansion of the minimal energy, in terms of the unknown (except in 1D) constants \( \min W \), \( \min \mathcal{W} \), and we have seen that – modulo the logarithmic terms in dimensions 1 and 2 – this next order term lives at the order \( n^{2-2/d} \), something which was not immediate. Moreover, with the statement “for \( P\text{-a.e. } (x,E) \), \( E \) minimizes \( \mathcal{W} \) over \( A_{\mu_0(x)} \),” we have obtained a characterization of the minimizers at the microscopic level: after blow-up around a point \( x \in \Sigma \) chosen uniformly, one sees a jellium of points with density (or “background”) \( \mu_0(x) \) (note that this is the only way – through the equilibrium measure – that the result depends on the potential \( V \)), which has to almost everywhere be a minimizer of \( \mathcal{W} \). This reduces the study of the minimizers of \( H_n \) to the minimization of \( \mathcal{W} \). If one believes that minimizers of \( \mathcal{W} \) are lattices, as one is led to believe in dimension 2 (recall in dimension 2 the Abrikosov triangular lattice is the unique best one), then one expects that minimizers of \( H_n \) microscopically look “almost-everywhere” like such lattices. Again, as in RNSe the “almost-everywhere” part of the statement should finally be replaceable by “everywhere.”

### 6.4 Consequences on the statistical mechanics

Now that we have completed the expansion at next order of the Hamiltonian \( H_n \), we may, just as in Section 2.6, apply it to obtain information on the thermal states and the partition function, by simply inserting this expansion into (2.58). In this context, the result of Theorem 6.4 on ground states can be viewed as a zero temperature (or \( \beta = \infty \)) result.

**Theorem 6.5 (Next order asymptotic expansion of the partition function).** Assume \( V \) is continuous and satisfies (A3)–(A6). Assume \( \bar{\beta} := \limsup_{n \to +\infty} \beta n^{1-2/d} > 0 \). Then there exists \( C_{\bar{\beta}} \) (depending on \( V \) and \( d \)) such that \( \lim_{\beta \to +\infty} C_{\bar{\beta}} = 0 \) and

\[
\log Z_{n,\beta} + \frac{\beta}{2} \left( n^2 I(\mu_0) - \left( \frac{n}{2} \log n \right) 1_{d=2} + n^{2-2/d} \xi_d \right) \leq C_{\bar{\beta}} \beta n, \tag{6.31}
\]

where \( \xi_d \) is as in (6.16).
In [SS8] one finds the analogous result in dimension 1: there exists $C_\beta$ with $\lim_{\beta \to +\infty} C_\beta = 0$ such that

$$|\log Z_{n,\beta} + \frac{\beta}{2} \left( n^2 I(\mu_0) - n \log n + n \xi_1 \right)| \leq C_\beta \beta n,$$

with $\xi_1 = \frac{1}{\beta^2} \min_{\mathcal{A}} W - \int_{\mathcal{B}} \mu_0 \log \mu_0$.

These results should be compared to the expansion stated in Theorem 2.3: again, we obtain here an expansion to next order. The expansion gets more precise as $\bar{\beta} \to +\infty$, i.e., $\beta \gg n^{2/d-1}$ (and then one essentially recovers the minimal energy as in Theorem 6.4). However, it already identifies the conjectured crystallization regime as $\beta \gg n^{2/d-1}$, and complements the result of Corollary 4.7 by a lower bound of similar form. Except in dimension 1, where as already mentioned, $Z_{n,\beta}$ is exactly known via Selberg integrals (at least for $V$ quadratic), these results improve on the known results.

We conclude with a result of large deviations type, which improves on Theorem 2.3 by bounding the probability of rare microscopic events. This is however not a complete Large Deviations Principle at next order, the search of which remains an open question.

**Theorem 6.6** (Rare events at the microscopic scale). Assume $V$ is continuous and satisfies (A3)–(A6). Let $i_n$ be the map which to any $x_1, \ldots, x_n \in \mathbb{R}^d$ associates $P_n \in \mathcal{P}(\Sigma \times X)$ as in Theorem 6.3. For any $n > 0$ let $A_n \subset (\mathbb{R}^d)^n$ and

$$A_\infty = \{ P \in \mathcal{P}(\Sigma \times X) | \exists (x_1, \ldots, x_n) \in A_n, i_n(x_1, \ldots, x_n) \to P \text{ up to a subsequence}\}.$$

Let $\bar{\beta} > 0$ be as in Theorem 6.5, $\xi_d$ be as in (6.16), and $\tilde{W}$ as in (6.14). There exists $C_\beta$ such that $\lim_{\beta \to +\infty} C_\beta = 0$, and

$$\limsup_{n \to \infty} \frac{\log P_\beta^A(A_n)}{n^{2/d-1}} \leq -\frac{\beta}{2} \left( \inf_{P \in A_\infty} \tilde{W} - \xi_d - C_\beta \right). \quad (6.32)$$

Here $A_\infty$ is the set of accumulation points of the $P_n$’s associated to configurations in $A_n$. Another way of phrasing this result is that in the limit $n \to \infty$

$$\tilde{W}(P) \leq \xi_d + C_\beta = \min \tilde{W} + C_\beta$$

except with exponentially decaying probability. This means that we have a threshold phenomenon: the Gibbs measure concentrates on configurations whose $\tilde{W}$ is below the minimum plus $C_\beta$. The threshold tends to 0 as $\beta \gg n^{2/d-1}$ and, in that regime, the Gibbs measure concentrates on minimizers of $\tilde{W}$, a weak crystallization statement. This indicates that even at nonzero temperature, configurations have some order (if one believes of seeing $\mathcal{W}$ as a measure of disorder, cf. the end of Section 5.4), in the sense that their $\tilde{W}$ cannot be too large.
Proof of Theorems 6.4 and 6.6. The proof is a direct consequence of the definition of \( \mathbb{P}^\alpha_n \) and Theorem 6.3, following the same method as in Theorem 2.3. We start by obtaining a lower bound for \( Z_{n,\beta} \) from Proposition 6.6: let \( \varepsilon > 0 \) be given. By definition of \( Z_{n,\beta} \), we may write

\[
Z_{n,\beta} = \int_{\mathbb{R}^d} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n \geq \int_{A_n} e^{-\frac{\beta}{2} H_n(x_1, \ldots, x_n)} dx_1 \ldots dx_n
\]

where \( A_n \) is the set given by Proposition 6.6. Taking the logarithm and inserting (6.25), we obtain

\[
\log Z_{n,\beta} \geq \log |A_n| - \frac{\beta}{2} \left( n^2 I(\mu_0) + \left( \frac{n}{2} \log n \right) 1_{d=2} + n^{2-2/d} (\xi_d + \varepsilon) \right).
\]

But \( \log |A_n| \geq \log(n!) r_1 d/n \geq -C\varepsilon n - C \) using Stirling’s formula, where \( C_\varepsilon \) depends on \( r_1 \), itself depending on \( \varepsilon \). Inserting, we obtain the lower bound corresponding to (6.31).

Conversely, let \( A_n \) be an arbitrary set in \( (\mathbb{R}^d)^n \). Assume \( (x_1, \ldots, x_n) \) minimizes (or almost minimizes) \( H_n \) over \( A_n \), then Theorem 6.3 gives us that

\[
H_n(x_1, \ldots, x_n) = 2n \sum_{i=1}^n \zeta(x_i) + \left( \frac{n}{2} \log n \right) 1_{d=2}
\]

\[
\geq n^{2-2/d} \left( \bar{W}(P) + o_n(1) \right) \geq n^{2-2/d} \left( \inf_{P \in A_\infty} \bar{W}(P) + o_n(1) \right).
\]

It follows that, inserting this estimate into (2.58), we obtain an upper bound on the probability of \( A_n \) by writing

\[
\log \mathbb{P}^\beta_n(A_n) \leq -\log Z_{n,\beta} - \frac{\beta}{2} \left( -\left( \frac{n}{2} \log n \right) 1_{d=2} + n^{2-2/d} \left( \inf_{P \in A_\infty} \bar{W}(P) + o_n(1) \right) \right)
\]

\[
+ \log \int_{\mathbb{R}^d} e^{-\beta n \sum_{i=1}^n \zeta(x_i)} dx_1 \ldots dx_n.
\]

Inserting the lower bound on \( \log Z_{n,\beta} \) obtained above, and the result of Lemma 4.8 (since our assumptions ensure that \( \beta n \to \infty \) as \( n \to \infty \)) we find

\[
\log \mathbb{P}^\beta_n(A_n) \leq -\log Z_{n,\beta} - \frac{\beta}{2} n^{2-2/d} \left( -\xi_d + \inf_{P \in A_\infty} \bar{W}(P) + o_n(1) \right)
\]

\[
+ n(|\omega| + o_n(1)). \quad (6.33)
\]

Taking in particular \( A_n = (\mathbb{R}^d)^n \), we have \( \mathbb{P}^\beta_n(A_n) = 1 \), and we can check that

\[
A_\infty \subset \left\{ P | E \in \bar{A}_{\mu_0(x)} P \text{-a.e, and the first marginal of } P \text{ is the normalized Lebesgue measure on } \Sigma \right\}
\]
by Theorem 6.3. It follows as previously that
\[
\inf_{P \in A_\infty} \tilde{W}(P) = \frac{1}{c_d} \int_{\mu_{\beta}(x)} \min A \, dx = \xi_d
\]
by the change of scales formula. Inserting into (6.33), we deduce the upper bound for $\log Z_{n,\beta}$ stated in (6.31). This completes the proof of (6.31) and of Theorem 6.4. Using then again (6.33) for a general $A_n$, and plugging in (6.31), we obtain the result of Theorem 6.6.

In order to obtain the exact asymptotic expansion of $\log Z_{n,\beta}$ at that order, one would need to understand how to compute the volume of the sets in configuration space (in terms of their electric currents). We expect this to create an entropy-type term. This is one of the main questions that remain open, and the object of future investigation [LS].
7 The Ginzburg-Landau functional: presentation and heuristics

In this chapter, we present some nonrigorous heuristics on the Ginzburg-Landau model that allow to see how and when vortices are expected to form in minimizers. A detailed presentation of the functional and of the related physics was already given in [SS4, Chap. 2], so we will try here to focus more on the new additions compared to that reference. We also refer to the classic book of [DeG].

7.1 The functional

Let us start by recalling the expression of the Ginzburg-Landau functional that was introduced in (1.1) in Chapter 1:

\[ G_\varepsilon(u, A) = \frac{1}{2} \int_\Omega |(\nabla - iA)u|^2 + |\text{curl } A - h_{\text{ex}}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2}, \]  

(7.1)

and the Ginzburg-Landau equations

\[ \begin{align*}
-\nabla^2 u &= \frac{1}{\varepsilon^2} u(1 - |u|^2) \quad \text{in } \Omega \\
-\nabla^\perp h &= \langle iu, \nabla u \rangle \quad \text{in } \Omega
\end{align*} \]  

(GL)

with boundary conditions

\[ \begin{align*}
\nabla A u \cdot \nu &= 0 \quad \text{on } \partial \Omega \\
h &= h_{\text{ex}} \quad \text{on } \partial \Omega.
\end{align*} \]

The precise meaning of the quantities appearing here was given in Chapter 1.

It is an important fact that this is a U(1) gauge theory, i.e. all the physically meaningful quantities, such as the energy and the equations (GL) are invariant under the gauge-transformations

\[ \begin{align*}
u \mapsto v e^{i\Phi} \\
A \mapsto A + \nabla \Phi
\end{align*} \]

where \( \Phi \) is any smooth real-valued function. One may easily check that, in addition to the energy, gauge-invariant quantities include \(|u|^2\) which represents the density of superconducting electrons, the magnetic field \( h = \nabla \times A \), and the superconducting current \( j := \langle iu, \nabla A u \rangle \). For more general gauge theories in theoretical physics, in particular non-Abelian ones, we refer to [JaTan] [MS].
7.2 Types of states, critical fields

7.2.1 Types of solutions and phase transitions

Three types of solutions (or states) to (GL) can be found:

1. the normal solution: \( (u \equiv 0, \text{curl} A \equiv h_{\text{ex}}) \). This is a true solution to (GL) and its energy is very easily computed: it is \( \frac{|\Omega|}{4\varepsilon^2} \).

2. the Meissner solution (or superconducting solution): \( (u \equiv 1, A \equiv 0) \), and all its gauge-transforms. This is a true solution if \( h_{\text{ex}} = 0 \), and a solution close to this one (i.e. with \( |u| \approx 1 \) everywhere) persists if \( h_{\text{ex}} \) is not too large. Its energy is approximately \( G_{\varepsilon}(1,0) = \frac{h_{\text{ex}}^2}{2} |\Omega| \). By comparing these energies, we see that the Meissner solution is more favorable when \( h_{\text{ex}} \) is small, while the normal solution is more favorable when \( h_{\text{ex}} \) is large enough, more precisely when \( h_{\text{ex}} > \frac{1}{\sqrt{2}} \).

3. the vortex solutions: there is another state, with vortices, called the mixed state where normal and superconducting phases co-exist, and which is more favorable for intermediate values of \( h_{\text{ex}} \).

The physics gives us the following more precise results. There are three main critical values of \( h_{\text{ex}} \) or critical fields \( H_{\text{c1}}, H_{\text{c2}}, \) and \( H_{\text{c3}} \), for which phase-transitions occur.

- For \( h_{\text{ex}} < H_{\text{c1}} \) there are no vortices and the energy minimizer is the superconducting state \( (u \equiv 1, A \equiv 0) \). (This is a true solution if \( h_{\text{ex}} = 0 \), and a solution close to this one (i.e. with \( |u| \approx 1 \) everywhere) persists if \( h_{\text{ex}} \) is not too large.) It is said that the superconductor “expels” the applied magnetic field, this is the “Meissner effect”, and the corresponding solution is called the Meissner solution.

- For \( h_{\text{ex}} = H_{\text{c1}} \), which is of the order of \( |\log \varepsilon| \) as \( \varepsilon \to 0 \), the first vortex(es) appear.

- For \( H_{\text{c1}} < h_{\text{ex}} < H_{\text{c2}} \) the superconductor is in the “mixed phase” i.e. there are vortices, surrounded by superconducting phase where \( |u| \approx 1 \). The higher \( h_{\text{ex}} > H_{\text{c1}} \), the more vortices there are. The vortices repel each other so they tend to arrange in triangular Abrikosov lattices in order to minimize their repulsion.

- For \( h_{\text{ex}} = H_{\text{c2}} \approx \frac{1}{\sqrt{\varepsilon}} \), the vortices are so densely packed that they overlap each other, and a second phase transition occurs, after which \( |u| \approx 0 \) inside the sample, i.e. all superconductivity in the bulk of the sample is lost.

- For \( H_{\text{c2}} < h_{\text{ex}} < H_{\text{c3}} \) superconductivity persists only near the boundary, this is called surface superconductivity. More details and the mathematical study of this transition are found in [FH] and references therein.

- For \( h_{\text{ex}} > H_{\text{c3}} = O(\frac{1}{\varepsilon^2}) \) (defined in decreasing fields), the sample is completely in the normal phase, corresponding to the “normal” solution \( u \equiv 0, h \equiv h_{\text{ex}} \) of (GL). See [GP] for a proof.

The picture below represents the phase diagram for two-dimensional superconductors which is found in physics textbooks.
7.2.2 Vortex solutions

We have already mentioned in Chapter 1 that vortices are isolated zeroes of the order parameter $u$, and that they come with an integer topological degree.

When $\varepsilon$ is small, the potential term in (7.1) implies that any discrepancy between $|u|$ and 1 is strongly penalized, and a scaling argument hints that $|u|$ is different from 1 only in regions of characteristic size $\varepsilon$. A typical vortex centered at a point $x_0$ “looks like” $u = \rho e^{i\varphi}$ with $\rho(x_0) = 0$ and $\rho(x) = f(\frac{|x-x_0|}{\varepsilon})$ where $f(0) = 0$ and $f$ tends to 1 as $r \to +\infty$, i.e. its characteristic core size is $\varepsilon$, and

$$\frac{1}{2\pi} \int_{\partial B(x_0, R\varepsilon)} \frac{\partial \varphi}{\partial r} = d \in \mathbb{Z} \quad (7.2)$$

is its degree (note that the phase $\varphi$ can only be understood as a “multi-valued function”). For example $\varphi = d\theta$ where $\theta$ is the polar angle centered at $x_0$ yields a vortex of degree $d$.

True radial solutions in $\mathbb{R}^2$ of the Ginzburg-Landau equations of degree $d$, of the form

$$u_d(r, \theta) = f_d(r)e^{id\theta}, \quad A_d(r, \theta) = g_d(r)(-\sin \theta, \cos \theta)$$

have been shown to exist \cite{Plo1,Plo2,BC}. In \cite{GS}, it was shown that they are stable for $\varepsilon < 1/\sqrt{2}$ and $|d| \leq 1$ and for $\varepsilon \geq 1/\sqrt{2}$, unstable otherwise.

Let us compute the approximate energy of the rescaled version of such a solution, in a bounded size domain (say $B_R$) : letting $\tilde{u}_d(r, \theta) = u_d(\frac{r}{\varepsilon}, \theta)$ and
\[ \widetilde{A}_d(r, \theta) = \frac{1}{d} A_d(\frac{r}{d}, \theta), \]

\[
\frac{1}{2} \int_{B_R} |\nabla \widetilde{A}_d \widetilde{u}_d|^2 = \frac{1}{2} \int_{B_{R/\varepsilon}} |\nabla A_d u_d|^2 \\
= \frac{1}{2} \int_{B_{R/\varepsilon}} |\nabla f_d|^2 + f_d^2 |\nabla (d\theta) - A_d|^2 \\
\approx \frac{1}{2} \int_{B_{R/\varepsilon}} |\nabla f_d|^2 + \frac{1}{2} \int_0^{R/\varepsilon} f_d^2 \left( \frac{d}{r} \right)^2 2\pi r dr d\theta \\
= \pi d^2 \int_1^{R/\varepsilon} \frac{dr}{r} = \pi d^2 \log \frac{R}{\varepsilon},
\]

as \( \varepsilon \to 0 \). Here we have used the fact that \( f_d \) is expected to have a cut-off effect in balls of lengthscale \( \varepsilon \) near the vortex center (here the origin). The error in the above estimate is in fact \( O(1) \) as \( \varepsilon \to 0 \). We thus see that with such an ansatz, in a solution with vortices, each of them “costs” at leading order an energy \( \pi d^2 |\log \varepsilon| \), with \( d \) its degree.

In a bounded domain, there are indeed solutions with several such vortices glued together, for example arranged along a triangular lattice (their existence is proved at least as a bifurcation from the normal solution, see [Chap2,Alm]).

7.2.3 Related models: superfluids and rotating Bose-Einstein condensates

For comparison, it is interesting to mention the energy functional corresponding to the Gross-Pitaevskii model (in the so-called Thomas-Fermi regime) of superfluids such as Bose-Einstein condensates, in rotation with velocity vector \( \Omega \), and in a confining potential \( V \) (cf. [Fe] for general reference and [LieSei2] for the derivation from quantum mechanics)

\[ \text{GP}(u) = \int_{\mathbb{R}^2} |\nabla u|^2 - (\Omega \times x) \cdot (iu, \nabla u) + V(x)|u|^2 + G|u|^4 \]

which, after completing the squares can equivalently be written as

\[ \text{GP}(u) = \int_{\mathbb{R}^2} |\nabla u - i\Omega \times xu|^2 + (V(x) - \Omega^2)|u|^2 + G|u|^4 \]  \hspace{1cm} (7.4)

or

\[ \text{GP}_\varepsilon(u) = \int_{\mathbb{R}^2} |\nabla u - i\Omega \times xu|^2 + \left( \frac{V_{\text{eff}}(x) - |u|^2}{2\varepsilon^2} \right)^2 \]  \hspace{1cm} (7.5)

for some effective potential \( V_{\text{eff}} \). The well-known analogy with the Ginzburg-Landau model is readily visible: the role of \( A \) is replaced by that of the angular momentum vector-field \( \Omega \times x \), and that of \( h_{\text{ex}} \) by \( \Omega \). The effective potential \( V_{\text{eff}} \) (which depends on \( V, \Omega, \) and \( G \)) does not create significant differences from the constant 1 when \( \varepsilon \) is small compared to other characteristic constants. The limit
\( \varepsilon \to 0 \) is called in this context the Thomas-Fermi regime. As a result of this strong analogy, the techniques developed for Ginzburg-Landau adapt well to treat such functionals, cf. the review article [CPRY] and references therein. The analogy functions well for rotation angles which are not too large, i.e. for the equivalent of the regime of \( h_{\text{ex}} \) (much) smaller than the second critical field \( H_{c2} \), but significantly breaks down after that, i.e. the physics is very different for very high rotation (but the mathematical tools are still useful).

### 7.3 Heuristics

#### 7.3.1 Rough heuristics

Let us examine the competition between all the (nonnegative) terms appearing in (7.1). We will write \( u \) in trigonometric form as \( u = \rho e^{i\phi} \), with again \( \phi \) a “multi-valued” function. For a configuration with vortices, we have the formal relation

\[
\nabla \times \nabla \phi = 2\pi \sum_i d_i \delta_{a_i},
\]

where the \( a_i \)'s are the vortex centers, and the \( d_i \in \mathbb{Z} \) their degrees. This relation is true in the sense of distributions. To check it, it suffices to test against a smooth function and use (7.2) (note also that \( \nabla \times \nabla = 0 \) for true functions).

As we have seen, the term \( \int (1-|u|^2)^2 \varepsilon^2 \) prefers Meissner states \( \rho \approx 1 \), or states with vortices of lengthscale scale \( \varepsilon \), while it disfavors the normal state \( \rho = 0 \). By explicit computation, the quantity \( |\nabla A u|^2 \) is in trigonometric form

\[
|\nabla A u|^2 = |\nabla \rho|^2 + \rho^2 |\nabla \phi - A|^2.
\]

The term \( \int_{\Omega} |\nabla A u|^2 \) thus favors \( \rho \) to be constant and

\[
\nabla \phi \approx A.
\]

The term \( \int |\nabla \times A - h_{\text{ex}}|^2 \) “prefers” the induced field \( h = \nabla \times A \) to “align” with the applied field \( h_{\text{ex}} \):

\[
\nabla \times A \approx h_{\text{ex}}
\]

Taking the curl of (7.8) and combining with (7.6) and (7.9) leads to the formal relation

\[
2\pi \sum_i d_i \delta_{a_i} \approx h_{\text{ex}},
\]

which indicates at least heuristically that when \( h_{\text{ex}} \) is not small, there should be vortices (otherwise the left-hand side would vanish). The question of understanding what (7.10) exactly means and in what sense, and what configurations of points \( a_i \) satisfy this assertion, is the core of the matter of our study. Based
on what we have seen in previous chapters, we can expect that the configuration of vortex points \( a_i \) (with degrees \( d_i = 1 \)) which best approximates the uniform distribution of density \( h_{\text{ex}} \) in (7.10), is the triangular Abrikosov lattice of density \( h_{\text{ex}} \). We will see that this becomes true only when \( h_{\text{ex}} \) is large enough, because as seen above, it costs a fixed amount \( \pi d_i^2 |\log \varepsilon| \) in the term \( \int |\nabla \varphi|^2 \) to create one vortex. This way, the size of \( A \), which is of order \( h_{\text{ex}} \) by (7.9), has to become significantly larger than \( |\log \varepsilon| \) for this heuristics to be completely correct. Below this threshold, boundary effects are also important, as we shall see, and the true optimal distribution of the vortices is a constant distribution in a subdomain of \( \Omega \), analogous to the equilibrium measure for Coulomb gases. We will next give more precise heuristics related to these facts, and give a complete proof in Chapter 9.

### 7.3.2 The vorticity measure and the London equation

The precise meaning to give to relations of the form (7.6) and (7.10) is given via the vorticity measure (or vorticity) of a configuration, defined by

\[
\mu(u, A) = \text{curl} \langle iu, \nabla_A u \rangle + h
\]

(7.11)

which is a gauge-invariant quantity. It was first introduced in this form in [SS2], and is the gauge-invariant analogue of the Jacobian determinant of \( u \) seen as a map from \( \mathbb{R}^2 \) to \( \mathbb{R}^2 \) in the outlook popularized by Jerrard and Soner [JS0], itself previously viewed in [BBH] as the Hopf differential of the map \( u \). It is also the analogue of the vorticity of a fluid. Note that in trigonometric representation we have

\[
\mu(u, A) = \text{curl} (j + A) = \text{curl} (\rho^2 (\nabla \varphi - A) + A) \approx \text{curl} \nabla \varphi
\]

at least when \( \rho \) is close to 1, as is expected in the limit \( \varepsilon \to 0 \). This is why, in view of (7.6), we may write the heuristic relation

\[
\mu(u, A) \approx 2\pi \sum_i d_i \delta_{a_i}.
\]

(7.12)

This relation holds in the asymptotics of \( \varepsilon \to 0 \), and its proper meaning will be given in the next chapter. Suffice it for now to say that it is more correct to replace the sum of Dirac masses in the right-hand side of (7.12) by Diracs smeared out at the scale \( \varepsilon \) – characteristic lengthscale of the vortices – that we denoted \( a^{(\varepsilon)}_i \), as alluded to in (1.3), and as done for the Coulomb gas.

Turning again to the functional (7.1), one may observe that for a fixed \( u \), the energy \( G_\varepsilon \) is a positive quadratic function of \( A \), thus always has a unique critical point in terms of \( A \), and that critical point is a minimum. We may thus always consider that without loss of generality, \( G_\varepsilon \) has been minimized with respect to \( A \), this decreases the energy and does not affect the zeroes of \( u \), i.e. the vortices and their degrees, which are the objects we wish to understand. This way, we may assume that the Euler-Lagrange equation associated with this minimization, which
is the second relation in (GL), is satisfied, together with its boundary condition:
\[
\begin{cases}
-\nabla \cdot h = j = \langle iu, \nabla_A u \rangle & \text{in } \Omega \\
\quad h = h_{\text{ex}} & \text{on } \partial \Omega.
\end{cases}
\] (7.13)

Taking the curl of this equation, one obtains \(-\Delta h = \text{curl } j\), which we can rewrite, by definition of \(\mu\), as
\[
\begin{cases}
-\Delta h + h = \mu(u, A) & \text{in } \Omega \\
\quad h = h_{\text{ex}} & \text{on } \partial \Omega.
\end{cases}
\] (7.14)

This is exactly the rigorous version of the London equation (1.3), which directly relates the vorticity \(\mu\) and the induced magnetic field \(h\). Another way of writing it is that
\[
h(x) = h_{\text{ex}} + \int_{\Omega} G_\Omega(x, y) (\mu(u, A) - h_{\text{ex}})(y) \, dy
\] (7.15)

where \(G_\Omega\) is a Green-type function (or more correctly a Yukawa potential) of the domain with Dirichlet boundary condition, solution to
\[
\begin{cases}
-\Delta G_\Omega(\cdot, y) + G_\Omega(\cdot, y) = \delta_y & \text{in } \Omega \\
G_\Omega(\cdot, y) = 0 & \text{on } \partial \Omega.
\end{cases}
\] (7.16)

This shows that \(h\) can be mathematically seen as the potential generated by the vorticity distribution \(\mu(u, A)\) via \(G_\Omega\). This kernel depends on the domain, but its leading contribution is the Coulomb kernel in dimension 2, \(-\frac{1}{2\pi} \log |x|\), hence the origin of the analogy with the 2D Coulomb gas, as explained in Chapter 1.

Note that when the vorticity vanishes, the London equation reduces to
\[
\begin{cases}
-\Delta h + h = 0 & \text{in } \Omega \\
\quad h = h_{\text{ex}} & \text{on } \partial \Omega,
\end{cases}
\] (7.17)

hence (up to dividing by \(h_{\text{ex}}\)) we can expect a particular role to be played by the function \(h_0\), solution of
\[
\begin{cases}
-\Delta h_0 + h_0 = 0 & \text{in } \Omega \\
\quad h_0 = 1 & \text{on } \partial \Omega,
\end{cases}
\] (7.18)

which depends only on the domain \(\Omega\) and exhibits exponential decay away from \(\partial \Omega\). The situation expected when there are no vortices is to have \(h \approx h_{\text{ex}} h_0\), which physically corresponds to the Meissner effect, for which it is said that the applied magnetic field is expelled by the superconducting sample and only penetrates in it in a layer localized near the boundary (in our normalization, this layer has characteristic lengthscale 1, but physically, it is the so-called penetration depth).

### 7.3.3 Approximation to the energy and formal derivation of the first critical field

We may now justify (1.2), which we recall here:
\[
G_\varepsilon(u, A) \approx \frac{1}{2} \int_{\Omega} |\nabla h|^2 + |h - h_{\text{ex}}|^2
\] (7.19)
with $h = \nabla \times A$ solving (7.14). Taking the norm of (7.13) we may compute that in trigonometric form $|\nabla h|^2 = |j|^2 = \rho^4 |\nabla \varphi - A|^2$. Comparing with (7.7) we thus have

$$\int_{\Omega} |\nabla h|^2 = \int_{\Omega} |\nabla \rho|^2 + \frac{|\nabla h|^2}{\rho^4}.$$  

But for any solution of (GL), it holds that $|u| \leq 1$ (this can be checked using the maximum principle on the equation satisfied by $|u|$). We may thus bound from below

$$\int_{\Omega} |\nabla h|^2 \geq \int_{\Omega} |\nabla \rho|^2$$

and we expect almost equality in view of the heuristic relation $\rho \approx 1$. The difference will turn out to be indeed negligible as $\varepsilon \to 0$ as a by-product of our analysis, based on comparing ansatz-free lower bounds and upper bounds obtained by explicit constructions.

Once (7.19) is established, deriving the first critical field can be done formally: at the point where the first vortices appear, we can expect that the induced magnetic field is well approximated to leading order by the magnetic field generated in the situation with no vortices, i.e. $h_{\text{ex}}h_0$. One may then expand around that function by setting $h = h_{\text{ex}}h_0 + h_{1,\varepsilon}$ where $h_{1,\varepsilon}$ is seen as a correction term, insert this into (7.19) and expand in terms of this splitting. This yields

$$G_{\varepsilon}(u_{\varepsilon}, A_{\varepsilon}) \approx \frac{h_{\text{ex}}^2}{2} \int_{\Omega} |\nabla h_0|^2 + |h_0 - 1|^2 + \frac{1}{2} \int_{\Omega} |\nabla h_{1,\varepsilon}|^2 + |h_{1,\varepsilon}|^2$$

$$+ h_{\text{ex}} \int_{\Omega} (-\Delta h_{1,\varepsilon} + h_{1,\varepsilon})(h_0 - 1)$$

$$= \frac{h_{\text{ex}}^2}{2} \int_{\Omega} |\nabla h_0|^2 + |h_0 - 1|^2 + \frac{1}{2} \int_{\Omega} |\nabla h_{1,\varepsilon}|^2 + |h_{1,\varepsilon}|^2$$

$$+ h_{\text{ex}} \int_{\Omega} (h_0 - 1)\mu(u_{\varepsilon}, A_{\varepsilon}),$$

where for the cross-term we have used an integration by parts, and (7.14) with (7.18). With the approximate relation (7.12) and estimating $\frac{1}{2} \int_{\Omega} |\nabla h_{1,\varepsilon}|^2 + h_{1,\varepsilon}^2$ as the cost to create a vortex, i.e. $\pi \sum_i d_i^2 |\log \varepsilon|$ by the heuristic of Section 7.2.2 we are led to

$$G_{\varepsilon}(u_{\varepsilon}, A_{\varepsilon}) \approx \frac{h_{\text{ex}}^2}{2} \int_{\Omega} |\nabla h_0|^2 + |h_0 - 1|^2 + \pi \sum_i d_i^2 |\log \varepsilon| + 2\pi h_{\text{ex}} \sum_i d_i (h_0 - 1)(a_i).$$

The energy of a configuration with vortices thus becomes smaller than that of the vortex-free Meissner solution if we can have

$$\pi \sum_i d_i^2 |\log \varepsilon| + 2\pi h_{\text{ex}} \sum_i d_i (h_0 - 1)(a_i) \leq 0.$$
Noting that $h_0 - 1 \leq 0$ in $\Omega$ by the maximum principle applied to the equation (7.18), a quick examination shows that this can be first achieved when

$$h_{\text{ex}} \geq \frac{|\log \varepsilon|}{2 \max_{\Omega} |h_0 - 1|}$$

and with vortices that have degrees $d_i = +1$, located at the point(s) where $h_0 - 1$ achieves its minimum (or equivalently $|h_0 - 1|$ achieves its maximum). This gives the leading order value of the first critical field

$$H_{c_1} \sim \frac{1}{2 \max_{\Omega} |h_0 - 1| |\log \varepsilon|} \quad \text{as } \varepsilon \to 0.$$  \hspace{1cm} (7.20)

This heuristic is in fact correct, it first appeared (in a slightly different but equivalent form) in [BR2], it was then justified rigorously in [Se1]. A proof of the most precise result can be found in [SS4, Chap. 12].

This expansion of $H_{c_1}$ confirmed and made more precise the expansion known to physicists, e.g. in [DeG], by giving the exact prefactor in terms of the domain $\Omega$, and locating the points of nucleation of the first vortices. As soon as there is more than one vortex accumulating near one of the optimal point(s), the Coulomb repulsion between vortices starts to kick in, and slightly delays the onset of more vortices. Again for details we refer to [SS4].

We will see in Chapter 9 how to derive more information about the number and optimal distribution of vortices above $H_{c_1}$.
8 Main mathematical tools for Ginzburg-Landau

Mathematicians started to get interested in the Ginzburg-Landau model mostly in the 90’s, with Berger-Chen, Chapman, Rubinstein, Schatzman, Du, Gunzburger, Baumann, Phillips, cf. e.g. to the review papers [Chap1, DGP]. Then, Bethuel-Brezis-Hélein [BBH] were the first to introduce tools to systematically study vortices, their exact profile and their asymptotic energy (with important input from Hervé-Hervé [HH] and Mironescu [Mi]). They did it in the simplified context of the two-dimensional Ginzburg-Landau equation not containing the magnetic gauge (set $A \equiv 0$, $h_{ex} \equiv 0$ in (GL)), and under an a priori bound $C|\log \varepsilon|$ on the energy, which allows only for a fixed number of vortices as $\varepsilon \to 0$. The analysis of that model was completed by many works, including the precise study of solutions by Comte-Mironescu [CM1, CM2], and the monograph of Pacard-Riviére [PaR].

The analysis of the simplified model was adapted to the model with gauge but with a different boundary condition and still the same a priori bound, by Bethuel and Riviére [BR1, BR2], still in dimension 2. The three-dimensional (more physical) versions were first studied in Riviére [Ri], and later in [LR, JSo, BBO, BJSO], among others. For a (slightly outdated) review, we refer to [SS4, Chap. 14].

An important challenge after these works was to be able to treat the case where the a priori bound is released and the number of vortices blows up as $\varepsilon \to 0$, as really happens in the full model with magnetic field. There are two main related technical tools that have been widely used and applied in such a situation. The first is the “vortex balls construction” method introduced independently by Jerard [Je] and Sandier [Sa], which allows to get completely general lower bounds for the energy of a configuration in terms of its vortices (regardless of their number and degrees). The second is the so-called “Jacobian estimate” which gives a quantitative estimate and meaning for (7.12), i.e. relates the vorticity of an arbitrary configuration, as defined in (7.11) (or the Jacobian in the gauge-independent version) to its underlying vortices.

8.1 The ball construction method

As mentioned, the ball construction method was first introduced in two slightly different variants in [Sa] and [Je], and it was reworked and improved over the years e.g. [SS8, SS4, SS5, JSn], and extended to higher dimensions [JSn, Sa2]. It would be too long here to prove the best-to-date results, but we will give an idea of the method and a statement of results. The main question in the end, for what we
need here, was to obtain estimates on the energy that allow for only an error of a constant per vortex.

8.1.1 A sketch of the method

The method consists in starting by understanding lower bounds for unit-valued complex functions in the plane.

A lower bound on an annulus Assume that \( u \) is a (complex-valued) function mapping an annulus (say centered at the origin) to the unit circle, in other words \( |u| = 1 \) in \( B_R \setminus B_r \). If \( u \) is sufficiently regular (say, continuous, for more refined assumptions, see \[BN1, BN2\] and references therein), we can define its degree as the integer

\[
d = \deg(u, \partial B_t) := \frac{1}{2\pi} \int_{\partial B_t} \langle iu, \frac{\partial u}{\partial \tau} \rangle = \frac{1}{2\pi} \int_{\partial B_t} \frac{\partial \varphi}{\partial \tau},
\]

which is constant over \( t \in [r, R] \) and where we have written \( u = e^{i\varphi} \) for \( \varphi \) some real-valued lifting of \( u \) (for questions of existence and regularity of a lifting see \[BM\] and references therein). We may then write

\[
\int_{B_R \setminus B_r} |\nabla u|^2 = \int_{B_R \setminus B_r} |\nabla \varphi|^2 \geq \int_r^R \int_{\partial B_t} \left| \frac{\partial \varphi}{\partial \tau} \right|^2 dt \quad (8.1)
\]

\[
\geq \int_r^R \left( \int_{\partial B_t} \frac{\partial \varphi}{\partial \tau} \right)^2 \frac{dt}{2\pi t} \quad (8.2)
\]

where the second relation follows by an application of Cauchy-Schwarz’s inequality. We then recognize the degree \( d \) and may write

\[
\int_{B_R \setminus B_r} |\nabla u|^2 \geq \int_r^R \frac{4\pi^2 d^2}{2\pi t} dt = 2\pi d^2 \log \frac{R}{r}, \quad (8.3)
\]

and there is equality if and only if \( \frac{\partial \varphi}{\partial \tau} \) is constant on each circle \( \partial B_t \), which amounts in the end to \( u \) being of the form \( e^{i(\theta + \theta_0)} \) in polar coordinates centered at the center of the annulus. The lower bound \( (8.3) \) is general and is the building block for the theory. It does show how a vortex of degree \( d \) induces a logarithmic cost, as in \( (7.3) \).

In a Ginzburg-Landau configuration with vortices, we will not have \( |u| = 1 \) everywhere, but we can expect that \( |u| \approx 1 \) except in small regions of scale \( \varepsilon \) around the vortex cores. We can expect to be able to localize the “bad regions” where \( |u| \) is far from 1, which must contain all the vortices, in balls of size \( C\varepsilon \). We may then center around each such ball an annulus of inner radius \( C\varepsilon \) and outer radius \( R \) (the largest possible so that it does not intersect any other vortex), and then the estimate \( (8.3) \) yields on such an annulus a lower bound by \( \pi d^2 \log \frac{R}{C\varepsilon} \sim \pi d^2 \log \frac{1}{\varepsilon} \) at leading order as \( \varepsilon \to 0 \). If we can build such annuli that are disjoint, then we
may add the lower bounds obtained this way and obtain a global lower bound of the form
\[ \int |\nabla u|^2 \geq \pi \sum_i d_i^2 (|\log |\varepsilon| + O(1)) \quad \text{as } \varepsilon \to 0 \]
where \(d_i\) are the degrees of the vortices. Two questions remain: first to find an algorithm to build such disjoint annuli in some optimal way, and second to handle the fact that we do not really have \(|u| = 1\) outside of small balls, but rather \(|u| \approx 1\), with a control via the energy term \(\int_{\Omega} \frac{(1-|u|^2)^2}{2\varepsilon^2}\). These questions are answered in a slightly different way by both the methods of [Je] and [Sa]; we now give the main elements.

**Construction of initial balls** To initiate the ball construction, one does need some weak upper bound on the energy, of the form
\[ G_\varepsilon(u, A) \leq C\varepsilon^\alpha - 1, \quad \alpha \in (0, 1) \]
which implies
\[ \int_{\Omega} |\nabla \rho|^2 + \frac{(1 - \rho^2)^2}{2\varepsilon^2} \leq C\varepsilon^\alpha \quad (8.4) \]
with \(\rho = |u|\). This control implies, for \(0 < \delta < 1\), a control of the total perimeter of the bad set \(\{\rho \leq 1 - \delta\}\) via the co-area formula (cf. [EvGar]), this is the argument used by Sandier: by Cauchy-Schwarz, we have
\[ \int_{\Omega} |\nabla \rho(1 - \rho^2) \leq C\varepsilon^\alpha \]
and the left-hand side is equal to
\[ \int_0^{+\infty} (1 - t^2) \mathcal{H}^1(\{|\rho(x)| = t\}) dt, \]
where \(\mathcal{H}^1\) denotes the one-dimensional Hausdorff measure. The upper bound thus allows to find many level sets \(\{\rho \leq 1 - \delta\}\), with \(\delta\) as small as a power of \(\varepsilon\), whose perimeter is small. A compact set of small perimeter can then be covered by disjoint closed balls \(B_i\) of radii \(r_i\), with \(\sum r_i\) controlled by that perimeter.

In Jerrard’s construction, the initial balls are obtained differently. The use of the co-area formula is replaced by the following lemma, based on elementary arguments:

**Lemma 8.1** (Lower bound on circles [Je], Lemma 2.3). Letting \(\rho\) be a real-valued function defined over \(B(x, r) \subset \mathbb{R}^2\) with \(r \geq \frac{1}{2}\varepsilon\), if \(m = \min_{\partial B(x, r)} \rho(x)\), we have
\[ \int_{\partial B(x, r)} |\nabla \rho|^2 + \frac{(1 - \rho^2)^2}{2\varepsilon^2} \geq c_0 \frac{(1 - m)^2}{\varepsilon} \]
for some universal constant \(c_0 > 0\).
This is another way of quantifying the cost of \(|u|\) being away from 1.

Then Jerrard only covers, again by disjoint closed balls of radii \(r_i\), the connected components of the set \(\{\rho \leq \frac{1}{2}\}\) on the boundary of which \(u\) has nonzero degree, and is able to do it in such a way that the radius \(r_i\) of each ball \(B_i\) is bounded above by \(\varepsilon\) times the energy that the ball contains.

**Ball construction method** Consider a collection of (disjoint closed) initial balls \(B_0 = \{B_i\}\) of radii \(r_i\), and let us assume to fix ideas that \(|u| = \rho = 1\) outside of these balls. If we have disjoint annuli centered around these same balls, of inner radii \(r_i\) and outer radii \(R_i\), then we may add the lower bounds given by (8.3) to obtain

\[
\frac{1}{2} \int_{\Omega} |\nabla u|^2 \geq \pi \sum_i d_i^2 \log \frac{R_i}{r_i},
\]

where \(d_i\) is the degree of \(u\) on each annulus. We then see that these lower bounds combine nicely if all the ratios \(\frac{R_i}{r_i}\) are equal, because then

\[
\log \frac{R_i}{r_i} = \log \frac{\sum_i R_i}{\sum_i r_i} = \log s
\]

where \(s\) is the common ratio \(R_i/r_i\), in other words the common conformal factor of the annuli. Let us underline that this is the point where the construction is purely two-dimensional: in higher dimensions the energy \(\int |\nabla u|^2\) is not conformally invariant and the estimates on annuli would not involve logarithms and not combine well.

In order to apply this reasoning, the annuli need to all be disjoint. The idea of the ball construction method is to build such disjoint annuli by continuously growing jointly all the initial balls, keeping their centers fixed, and multiplying their radii by the same factor \(s \geq 1\), until \(s\) is large enough (typically of order \(1/\varepsilon\)). This way the previous reasoning applies, and at least for \(s\) close enough to 1, the balls (hence the annuli) remain disjoint.

At some point during the growth process, two (or more) balls can become tangent to each other. The method then is to merge them into a ball that contains them both and has a radius equal to the sum of the radii of the merged ball. In other words, if \(B_1 = B(a_1, r_1)\) and \(B_2 = B(a_2, r_2)\) are tangent, we merge them into \(B' = B\left(\frac{a_1 r_1 + a_2 r_2}{r_1 + r_2}, r_1 + r_2\right)\). (The resulting ball can then intersect other balls, in which case one proceeds to another merging, etc, until all the balls are disjoint). The merging process preserves the total sum of the radii, and as for the degrees we have

\[
\deg(u, \partial B') = \deg(u, \partial B_1) + \deg(u, \partial B_2).
\]  

Thus the only problem is that the \(d_i^2\) do not add up nicely during merging. The price to pay is to give up on obtaining a lower bound with a \(\sum_i d_i^2\) factor, but rather to keep a lower bound by the smaller factor \(\sum_i |d_i|\) (here we use that the \(d_i\)’s are all integers). Such factors do add up nicely through merging since we have (8.5) during a merging, thus \(|d| = |d_1 + d_2| \leq |d_1| + |d_2|\).
8.1 The ball construction method

The fact that we need to abandon the hope of lower bounds by $\sum d_i^2$ is completely natural, due to possible cancellations of singularities (or vortices) of $u$. If $u$ has a vortex of degree $+1$, and a nearby vortex of degree $-1$ at distance $r$, once the associated balls have been merged, the total degree is $0$, and one does not expect any substantial energy to lie in the annuli surrounding the merged balls.

After mergings, the old collection is replaced by the new collection (with merged balls) which is still made of disjoint balls, and the growth process is resumed, until some next intersection and merging happens, etc. The construction can then be stopped at any value of the parameter $s$, depending on the desired final total radius of the balls, and the desired final lower bound.

Using this method and combining it with (8.3), one arrives at the following result, where for any ball $B$, $r(B)$ will denote its radius. Also if $B$ is a collection of balls, $\lambda B$ is the collection of balls with same centers, and radii multiplied by $\lambda$.

**Proposition 8.2** (Ball construction). Let $B_0$ be a collection of disjoint closed balls in the plane. Assume $u : \Omega \setminus \bigcup_{B \in B_0} \rightarrow S^1$. For any $s \geq 1$ there exists a family of balls $B(s)$ such that the following holds.

- $B(1) = B_0$.
- For any $s_1 \leq s_2$ we have
  \[ \bigcup_{B \in B(s_1)} B \subset \bigcup_{B \in B(s_2)} B. \]
- There exists a finite set $T$ (the set of merging “times”) such that if $[s_1, s_2] \subset [1, +\infty) \setminus T$, we have $B(s_2) = \frac{s_2}{s_1} B(s_1)$.
- \[ \sum_{B \in B(s)} r(B) = s \sum_{B \in B_0} r(B). \]
- For any $B \in B(s)$ such that $B \subset \Omega$, denoting $d_B = \deg(u, \partial B)$ we have
  \[ \frac{1}{2} \int_B |\nabla u|^2 \geq \pi |d_B| \log s = \pi |d_B| \log \frac{\sum_{B \in B(s)} r(B)}{\sum_{B \in B_0} r(B)}. \]

This is the building block estimate. As mentioned, one needs to control the initial total radius by some small factor; typically, one can expect it to be $n \epsilon$ where $n$ is the number of initial balls. Then one may choose the parameter $s$ according to the needs, so that the final sum of the radii be not too large, but still large enough for the factor in the right-hand side to be at leading order $\pi |d| \log \frac{1}{2\epsilon}$. For example, a good choice may be $\sum_{B \in B(s)} r(B) = 1/|\log \epsilon|$, which is a $o(1)$ quantity (guaranteeing small balls) but such that $\log(\sum_B r(B))^{-1} = \log |\log \epsilon|$ is negligible compared to $|\log \epsilon|$. 

Dealing with non unit-valued functions The main technical difficulty that remains is to handle the fact that $|u|$ is not really equal to 1 outside of small “initial balls.” One then needs to use the fact that outside of the initial balls $u$ does not vanish and one may write

$$
\int |\nabla u|^2 = \int \rho^2 \left| \nabla \left( \frac{u}{\rho} \right) \right|^2 + |\nabla \rho|^2 + \frac{(1-\rho^2)^2}{2\varepsilon^2}, \quad (8.6)
$$

use a bound from below for $\rho$, and then bounds from below for unit vector fields to bound from below $\int |\nabla (\frac{u}{\rho})|^2$.

In Sandier’s construction, this is handled by combining the result of Proposition 8.2 with a co-area argument as outlined above, but in a rather sophisticated manner since the argument has to be applied to all sub-level sets at once. For details, we refer to [SS4, Chap. 4].

In Jerrard’s construction, this is handled by combining (8.6) with the result of Lemma 8.1 and (8.3) to obtain

$$
\int_{\partial B_r} |\nabla u|^2 + \frac{(1-|u|^2)^2}{2\varepsilon^2} \geq m^2 \frac{2\pi d^2}{r} + c_0 \frac{(1-m)^2}{\varepsilon},
$$

with $m = \min_{\partial B_r} |u|$. Optimizing over $m$ yields

$$
\frac{1}{2} \int_{\partial B_r} |\nabla u|^2 + \frac{(1-|u|^2)^2}{2\varepsilon^2} \geq \lambda_\varepsilon \left( \frac{r}{|d|} \right)
$$

with $\lambda_\varepsilon(s)$ that behaves like $\min(\frac{s}{\varepsilon}, \frac{s}{\pi})$, and whose antiderivative $\Lambda_\varepsilon$ satisfies $\Lambda_\varepsilon(s) \geq \pi \log \frac{s}{\varepsilon} - C$. The balls are grown and merged (in the same way as explained before) in such a way that it’s not the factor of sum of radii which is constant, but rather the parameter $s = r(B)/|d_B|$, to preserve $r \geq s|d|$. One checks that the estimate

$$
\frac{1}{2} \int_B |\nabla u|^2 + \frac{(1-|u|^2)^2}{2\varepsilon^2} \geq r(B) \frac{\Lambda_\varepsilon(s)}{s} \geq \pi |d| \left( \frac{s}{\varepsilon} - C \right)
$$

is true initially and is preserved through the growth and merging process, yielding the desired estimate at the end of the growth process. For more details, we refer to [Je].

In all cases, the presence of the gauge $A$ does not change substantially the situation, the method consists in controlling the error that it creates via the term $\int |\text{curl} A - h_{\text{ex}}|^2$.

8.1.2 A final statement

To give a more precise idea, let us now finish with the statements of a result on the complete Ginzburg-Landau functional. If one is interested in the Ginzburg-Landau functional without magnetic gauge, then it suffices to set $A \equiv 0$ in the following result.
8.1 The ball construction method

It is borrowed from [SS4, Theorem 4.1]. A similar result (slightly stronger in some sense, slightly weaker in some other) and following Jerrard’s construction [Je], can be found in [SS5, Proposition 5.2]. We let $F_\varepsilon$ denote the Ginzburg-Landau energy with $h_{\text{ex}}$ set to 0.

**Theorem 8.1** (Ball construction lower bound). For any $\alpha \in (0,1)$, there exists $\varepsilon_0(\alpha)$ such that for any $\varepsilon < \varepsilon_0$, if $(u, A)$ is such that $\int_{\Omega} |\nabla u|^2 + \frac{(1-|u|^2)^2}{2\varepsilon^2} \leq \varepsilon^{\alpha - 1}$, the following holds.

For any $1 > r > \varepsilon^{\alpha / 2}$, there exists a finite collection of disjoint closed balls $B = \{B_i\}$ such that

- $\sum_{B \in B} r(B) = r$
- $\{x \in \Omega | \text{dist}(x, \partial \Omega) > \varepsilon, |u(x)| - 1| \geq \varepsilon^{\alpha / 4}\} \subset \bigcup_i B_i$.

- Writing $d_i = \text{deg}(u, \partial B_i)$ if $B_i \subset \{x \in \Omega | \text{dist}(x, \partial \Omega) > \varepsilon\}$, $d_i = 0$ otherwise, and $D = \sum_i |d_i|$, we have
  \[
  \frac{1}{2} \int_{\bigcup_i B_i} |\nabla A u|^2 + |\text{curl} A|^2 + \frac{(1-|u|^2)^2}{2\varepsilon^2} \geq \pi D \left( \log \frac{r}{D\varepsilon} - C \right) \tag{8.7}
  \]
  where $C$ is a universal constant.
- If in addition $F_\varepsilon(u, A) \leq \varepsilon^{\alpha - 1}$ then
  \[
  D \leq C \frac{F_\varepsilon(u, A)}{\alpha |\log \varepsilon|}.
  \]

In practice the last item already gives a rough lower bound on the energy $F_\varepsilon$ (without optimal constants) which can serve to provide a first control on $D$, which can then be inserted into the main result (8.7). Compared to the heuristic lower bound of $\pi \sum_i d_i^2 |\log \varepsilon|$, this lower bound

- loses $d_i^2$ and replaces it by $|d_i|$; as explained this is normal due to possible cancellations between vortices happening at small scales.
- introduces an error $-D \log D$: this is also normal due to the possibility of many vortices accumulating near a point, or near the boundary (think of the case of $n$ vortices of degree 1 regularly placed at small distance from the the boundary of the domain).
- introduces an error $-CD$ where $C$ is an unknown constant. When one knows that the number of vortices is bounded, the analysis derived from [BBH] allows to identify the constant order term in the energy of a vortex. It is (at least in the case of degree $\pm 1$ vortices), a constant that they denote $\gamma$, and which depends on the explicit optimal profile of the modulus of $u$ for a radial vortex (identified in [HH][MI]). One thus usually proceeds in two steps: first control the number of vortices via ball construction lower bounds which give the correct leading order energy, then if one can show that the number of vortices is locally bounded, recover this constant order term $\gamma$. 


• In order to accomplish this program, one may need (we needed) to localize the above lower bound over finite size balls in a (possible large) domain, and eliminate the $-D \log D$ error. We have seen that the energy carried by vortices is not only located at the vortex centers, it is spread over relatively large annuli surrounding them. In case of vortices accumulating near a point, the ball construction (because it stops at finite total radius) is missing some energy $\pi D^2 \log \frac{R}{r}$ (as given by (8.3)) which is lying in even larger annular regions. Although $\log \frac{R}{r}$ is then of order 1, such an energy suffices to compensate $-D \log D$ thanks to the power 2 in $D^2$ which beats $-D \log D$ when $D$ gets large. The method to do this and combine it with the ball construction, itself properly localized, is quite technical in its details, and is the object of [SS5], to which we refer the interested reader.

8.2 The “Jacobian estimate”

Let us now turn to the “Jacobian estimate,” which allows to give a rigorous meaning to (7.12), in terms of the result of a ball construction. Estimates of the same nature already appeared in [BR1, Se1, SS2], the estimate was optimized and its name popularized by the work of Jerrard and Soner [JSo]. Let us state it in the version presented in [SS4, Chap. 6]

The case without gauge $A$ is again contained in what follows by taking $A \equiv 0$.

Theorem 8.2 (Jacobian estimate). Let $u : \Omega \to \mathbb{C}$ and $A : \Omega \to \mathbb{R}^2$ be $C^1$. Let $B$ be a collection of disjoint closed balls with centers $a_i$ and radii $r_i$ such that

$$\{ x \in \Omega, \text{dist}(x, \partial \Omega) > \varepsilon, ||u|| - 1 \geq \frac{1}{2} \} \subset \bigcup_{B \in B} B.$$

Then, letting $d_i = \text{deg}(u, \partial B(a_i, r_i))$ if $B(a_i, r_i) \subset \{ x \in \Omega, \text{dist}(x, \partial \Omega) > \varepsilon \}$ and 0 otherwise, defining $\mu(u, A)$ by (7.11), if $\varepsilon$ and $r$ are less than 1, we have for $C > 0$ some universal constant

$$\left\| \mu(u, A) - 2\pi \sum_i d_i \delta_{a_i} \right\|_{(C^0_0(\Omega))^*} \leq C \max(\varepsilon, \sum_i r_i)(1 + F_{\varepsilon}(u, A)). \quad (8.8)$$

Moreover $\|\mu(u, A)\|_{(C^0_0(\Omega))^*} \leq CF_{\varepsilon}(u, A)$.

The spaces $(C^0)^*$ and $(C^{0,1})^*$ here are the space of bounded Radon measures and the dual of Lipschitz functions, respectively. For estimates in the dual of Hölder spaces, see the statement in [SS4]. Note that this result is naturally meant to work with a collection of disjoint balls obtained by a ball-construction. The total radius chosen to end the construction has to be taken small enough if one wants the estimate to be precise — this is of course in competition with the lower
bound estimate which improves as the total radius gets larger. So we see why
\( \sum_i r_i \) has to be optimized according to the needs. Note that it’s the centers of the
final balls in the construction (that may depend on the final total radius chosen \( r \))
which play the role of approximate vortex centers. More precise estimates can be
obtained, but with points \( a_i \) that do not correspond to centers of balls obtained
in a ball construction, this was done in [JS].

The proof is easy enough that we can give its main argument.

**Proof.** We set \( \mu = \mu(u, A) \). First, let us consider the function \( \chi \) on \([0, +\infty]\) defined
by \( \chi(t) = 2t \) if \( t \leq 1/2 \), \( \chi(t) = 1 \) if \( t \in [1/2, 3/4] \), \( \chi(t) = t \) if \( t \geq 2 \), and \( \chi \) is continuous
and piecewise affine. It satisfies

\[
|\chi(t)|^2 - t^2| \leq 3t|1-t|. \tag{8.9}
\]

We may then set \( \tilde{\mu} = \chi(|u|)\frac{\mu}{|\mu|} \). By assumption on the balls, we have \( |\tilde{u}| = 1 \)
outside of \( \cup_{B \in B} B \). We then define \( \tilde{\mu} = \text{curl} \langle i\tilde{u}, \nabla A\tilde{u} \rangle + \text{curl} A \) and check two
facts :

\[
\|\mu - \tilde{\mu}\|_{C^{\alpha,1}(\Omega)} \leq C\varepsilon F_{\varepsilon}(u, A), \tag{8.10}
\]

\[
\tilde{\mu} = 0 \text{ outside } \cup_{B \in B} B. \tag{8.11}
\]

For the first fact, it suffices to use an integration by parts: let \( \zeta \) be a smooth
test-function vanishing on \( \partial \Omega \). By definition of \( \tilde{u} \) we have

\[
\left| \int_{\Omega} \zeta (\mu - \tilde{\mu}) \right| = \left| \int_{\Omega} \nabla^{\perp} \cdot (\langle iu, \nabla A u \rangle - \langle i\tilde{u}, \nabla A \tilde{u} \rangle) \right|
\leq \|\nabla \zeta\|_{L^\infty(\Omega)} \int_{\Omega} \frac{|u|^2 - |\tilde{u}|^2}{|u|} |\nabla A u| \leq 3\|\nabla \zeta\|_{L^\infty(\Omega)} \int_{\Omega} |1 - |u|| |\nabla A u| \left| \tag{8.12}
\right.
\]

where we used the formal relation \( \langle iu, \nabla A u \rangle = \rho^2(\nabla \varphi - A) \) and \( |\nabla A u| \geq \rho |\nabla \varphi - A| \)
together with \( \text{[8.0]} \). It then suffices to apply Cauchy-Schwarz to control the right-hand
side of \( \text{[8.12]} \) by \( \varepsilon F_{\varepsilon}(u, A) \) and conclude \( \text{[8.10]} \). \( \text{[8.11]} \) is a consequence of
the simple observation that wherever \( |u| = 1 \), we have \( \text{curl} \langle iu, \nabla A u \rangle + \text{curl} A = \text{curl} \nabla \varphi = 0 \). We thus know that \( \tilde{\mu} \) is supported in the (disjoint) balls only, and
thus we may write, for any smooth test-function \( \zeta \) vanishing on \( \partial \Omega \),

\[
\int_{\Omega} \zeta \tilde{\mu} = \sum_i \int_{B(a_i, r_i)} \zeta \tilde{\mu} = \sum_i \zeta(a_i) \int_{B(a_i, r_i)} \text{curl} \langle i\tilde{u}, \nabla A \tilde{u} \rangle + \sum_i \int_{B(a_i, r_i)} (\zeta - \zeta(a_i))\tilde{\mu}.
\]

The first term of the right-hand side of this relation can be handled by Stokes’
theorem, and recalling that \( |\tilde{u}| = 1 \) on the boundary of each ball and the definition of
the degree, we find

\[
\sum_i \zeta(a_i) \int_{B(a_i, r_i)} \text{curl} \langle i\tilde{u}, \nabla A \tilde{u} \rangle + A = 2\pi \sum_i d_i \zeta(a_i) \tag{8.13}
\]
(for the balls that are $\varepsilon$-close to $\partial \Omega$ we need to replace $a_i$ by the nearest point on the boundary). The second term can be bounded above thanks to the Lipschitz continuity of $\zeta$ by
\[
\|\zeta\|_{C^{0,1}(\Omega)} \sum_i r_i \int_{B(a_i, r_i)} |\tilde{\mu}|.
\]
Noting that $\tilde{\mu} = 2(\partial_x \tilde{u} - iA_x \tilde{u}) \times (\partial_y \tilde{u} - iA_y \tilde{u}) + \text{curl } A$ (this is the same as using the formal relation $\text{curl } \langle i\tilde{u}, \nabla \tilde{u} \rangle = \text{curl } \langle \tilde{\rho}^2(\nabla \varphi - A) \rangle = \nabla \times \tilde{\rho}^2 \cdot \nabla \varphi$), we can bound $|\tilde{\mu}|$ by $|\nabla A \tilde{u}|^2 + |h|$, and in view of the definition of $\chi$ we are led to the control of the second term by $\|\zeta\|_{C^{0,1}(\Omega)} (\sum_i r_i + \varepsilon) (1 + F_{\chi}(u, A))$ (a little discussion is again needed for the balls that are very close to the boundary). Combining this with (8.10) and (8.13), we obtain the result.
9 The leading order behavior for Ginzburg-Landau

In this chapter, thanks to the tools presented in the previous chapter, we carry out the same program as in Chapter 2 i.e. the program of obtaining a mean-field limit or leading order behavior of minimizers (or ground states) of the Ginzburg-Landau functional (without temperature). The content of this chapter is essentially that of [SS2] or [SS4, Chap. 7], but we will try here to highlight the analogy with the Coulomb gas.

9.1 The $\Gamma$-convergence result

In what follows, the space $H^{-1}(\Omega)$ denotes the dual of the Sobolev space $H_0^1(\Omega)$, and $\mathcal{M}(\Omega)$ denotes the space of bounded Radon measures over $\Omega$, i.e. $C^0(\Omega)^*$. For a measure $\mu$ in $\mathcal{M}(\Omega)$, $|\mu|(\Omega)$ denotes its total variation.

We admit the fact (see e.g. [SS4, Sec 7.3.1] or [GMS1, Lemma 3.2]) that if $\mu \in H^{-1}(\Omega)$ then $U_\mu(x) = \int G_\Omega(x,y) \, d\mu(y)$, with $G_\Omega$ given by (7.16) makes sense and we have

$$\int_\Omega |\nabla U_\mu|^2 + |U_\mu|^2 = \int_{\Omega \times \Omega} G_\Omega(x,y) \, d\mu(x) \, d\mu(y). \tag{9.1}$$

**Theorem 9.1** ($\Gamma$-convergence of the Ginzburg-Landau functional [SS2], [SS4], Chap. 7). Assume

$$\lim_{\varepsilon \to 0} \frac{h_{\text{ex}} |\log \varepsilon|}{\varepsilon} = \lambda > 0. \tag{9.2}$$

Then, as $\varepsilon \to 0$, the functional $\frac{G_{\varepsilon}}{h_{\text{ex}}}$ $\Gamma$-converges as $\varepsilon \to 0$, for the sense of the convergence of $(u, A_\varepsilon)$ to $u$ in $\mathcal{M}(\Omega)$, to the functional

$$E_\lambda(\mu) = \frac{1}{2\lambda} |\mu|(\Omega) + \frac{1}{2} \int_\Omega |\nabla h_\mu|^2 + |h_\mu - 1|^2 \tag{9.3}$$

defined over $\mathcal{M}(\Omega) \cap H^{-1}(\Omega)$, where $h_\mu$ is the potential generated by $\mu$ as follows:

$$\begin{cases} -\Delta h_\mu + h_\mu = \mu & \text{in } \Omega \\ h_\mu = 1 & \text{on } \partial \Omega. \end{cases} \tag{9.4}$$

**Remark 9.1.** 1. Note that here we use a sense of convergence of $(u, A)$ that is the convergence of a nonlinear function of $(u, A)$, as in Remark 2.5 in Chapter 2. This is otherwise the counterpart of Proposition 2.8 for Ginzburg-Landau.
2. We could obtain convergence in a stronger sense for \( \mu(u_\varepsilon, A_\varepsilon)/h_{\text{ex}} \), we refer to [SS4, Chap. 7].

3. We can in fact obtain the same result when \( \lambda = \infty \), provided \( h_{\text{ex}} \ll \frac{1}{\varepsilon^2} \) as \( \varepsilon \to 0 \). This is done in [SS4, Chap. 8].

9.2 The proof of \( \Gamma \)-convergence

As in every proof of \( \Gamma \)-convergence, we need to prove a lower bound and an upper bound through the construction of a recovery sequence.

9.2.1 Lower bound

We in fact prove the stronger result of \( \Gamma \)-liminf + compactness (cf. Remark 2.2 in Chapter 2) i.e. that if \( \frac{1}{h_{\text{ex}}} G_\varepsilon(u_\varepsilon, A_\varepsilon) \) is bounded, then \( \mu(u_\varepsilon, A_\varepsilon)/h_{\text{ex}} \) has a convergent subsequence and the \( \Gamma \)-liminf relation holds. The lower bound relies on the estimates given by the ball construction method, and some lower semi-continuity arguments. Compared to the situation of the Coulomb gas, we do not have to worry about removing the diagonal terms, these are naturally smoothed out (at the scale \( \varepsilon \)) in the Ginzburg-Landau functional, but in turn we have to estimate these terms, corresponding to the self-interaction – or cost – of each vortex, and we do so via the ball construction method. Also we do not have an energy in the form of a sum of pairwise interactions but rather in integral form, as an integral of the potential generated by the charges, equivalent to (3.8). Note also that the fact that the vortex degrees do not have fixed sign would create difficulties in using the same method as in the proof of Proposition 2.8.

Let us start from an arbitrary family of configurations \( (u_\varepsilon, A_\varepsilon) \), assuming \( G_\varepsilon(u_\varepsilon, A_\varepsilon) \leq C \) for some \( C \) independent of \( \varepsilon \). Since we assume \( h_{\text{ex}} \sim \lambda |\log \varepsilon| \) with \( \lambda > 0 \), this also implies that \( h_{\text{ex}} \leq C |\log \varepsilon| \) and thus \( G_\varepsilon(u_\varepsilon, A_\varepsilon) \leq C |\log \varepsilon|^2 \).

We may then apply Theorem 8.1 with \( \alpha = \frac{1}{2} \) and final radius \( r = |\log \varepsilon|^{-10} \). It yields a collection of balls \( \{ B_i \} \) covering \( \Omega_{\varepsilon} := \{ x \in \Omega, \text{dist}(x, \partial \Omega) > \varepsilon \} \), outside of which we have \( |u_\varepsilon| - 1 \leq \varepsilon^{1/4} \), with \( \int |CF_\varepsilon(u, A)| \geq D |\log \varepsilon| \), and

\[
\int_{\bigcup_i B_i} |\nabla A_\varepsilon u_\varepsilon|^2 + |h_\varepsilon|^2 + \frac{(1 - |u_\varepsilon|^2)^2}{2\varepsilon^2} \geq \pi \sum_i |d_i| \left( \log \frac{1}{\sum_i |d_i|} - C \log |\log \varepsilon| \right).
\] (9.5)

One immediately checks that the bound \( D |\log \varepsilon| \leq CF_\varepsilon(u, A) \), the bounds on \( G_\varepsilon \) and on \( h_{\text{ex}} \) yield \( D \leq C |\log \varepsilon| \) for some constant \( C \) (depending only on \( \lambda \)). Plugging this into \( (9.5) \), we get

\[
\int_{\bigcup_i B_i} |\nabla A_\varepsilon u_\varepsilon|^2 + |h_\varepsilon - h_{\text{ex}}|^2 + \frac{(1 - |u_\varepsilon|^2)^2}{2\varepsilon^2} \geq \pi \sum_i |d_i| (|\log \varepsilon| - C \log |\log \varepsilon|).
\] (9.6)

\(^1\text{recall that } F_\varepsilon \text{ is the Ginzburg-Landau functional } G_\varepsilon \text{ with } h_{\text{ex}} \text{ set to 0} \)
It also implies that, defining \( \nu_{\varepsilon} = 2\pi \sum_i d_i \delta_{a_i} \) (the discrete approximate Jacobian), we have that \( \frac{\nu_{\varepsilon}}{h_{\varepsilon}} \) is bounded in the sense of measures (since (9.2) holds). Thus, up to extraction, we can assume that \( \frac{\nu_{\varepsilon}}{h_{\varepsilon}} \rightarrow \mu \) for some bounded Radon measure \( \mu \in \mathcal{M}(\Omega) \), and we have

$$\liminf_{\varepsilon \to 0} \frac{2\pi \sum_i |d_i|}{h_{\varepsilon}} \geq |\mu|(\Omega).$$  \tag{9.7}$$

In addition, with the Jacobian estimate Theorem 8.2 by choice of the final radius \( r = |\log \varepsilon|^{-1/4} \), we find that \( \frac{1}{h_{\varepsilon}} (\mu(u_{\varepsilon}, A_{\varepsilon}) - \nu_{\varepsilon}) \rightarrow 0 \) in \( (C^0,1(\Omega))^* \) and thus we also have

$$\frac{\mu(u_{\varepsilon}, A_{\varepsilon})}{h_{\varepsilon}} \rightarrow \mu \quad \text{in} \quad (C^0,1(\Omega))^*. \tag{9.8}$$

Next, since we are looking for a lower bound for \( G_{\varepsilon}(u, A) \) we can assume without loss of generality that \( G_{\varepsilon}(u, \cdot) \) has been minimized with respect to \( A \), which ensures, as explained in Chapter 7, that the second Ginzburg-Landau equation (7.13) is satisfied, hence also the London equation (7.14). Dividing (7.14) by \( h_{\varepsilon} \) and using (9.8), we find that \( \frac{h_{\varepsilon}}{h_{\varepsilon}} \) (where \( h_{\varepsilon} = \text{curl } A_{\varepsilon} \)) converges (say in the sense of distributions) to some \( \mu_{\varepsilon} \) which is related to \( \mu \) via (9.4).

We recall that (7.13) implies that \( |u_{\varepsilon}|^2 |\nabla A_{\varepsilon, u_{\varepsilon}}|^2 \geq |\nabla h_{\varepsilon}|^2 \). Since \( |u_{\varepsilon}| = 1 \) outside of the balls modulo an error \( \varepsilon^{1/4} \), we may bound from below \( \int_{\Omega \cup \cup B_1} |\nabla A_{\varepsilon, u_{\varepsilon}}|^2 \) by \( \int_{\Omega \cup \cup B_1} |\nabla h_{\varepsilon}|^2 + o(1) \), and thus with (9.6) we are led to

$$G_{\varepsilon}(u_{\varepsilon}, A_{\varepsilon}) \geq \pi \sum_i |d_i| (|\log \varepsilon| - C \log |\log \varepsilon|) + \int_{\Omega \cup \cup B_1} |\nabla h_{\varepsilon}|^2 + |h_{\varepsilon} - h_{\varepsilon}|^2.$$

The last step is to divide by \( h_{\varepsilon}^2 \) and pass to the liminf. For the second term, we observe that since \( \sum_i r_i \rightarrow 0 \) as \( \varepsilon \), we may extract a sequence \{\( \varepsilon_n \}\} such that \( \mathcal{A}_N := \cup_{n \geq N} (\cup B_i) \) satisfies \( |\mathcal{A}_N| \rightarrow 0 \) as \( N \rightarrow \infty \). In other words, there exists an arbitrarily small set which contains all the balls for all \( \varepsilon \)'s along the subsequence. We may then write

$$C \geq \liminf_{n \to \infty} \frac{G_{\varepsilon_n}(u_{\varepsilon_n}, A_{\varepsilon_n})}{h_{\varepsilon_n}^2} \geq \liminf_{\varepsilon \to 0} \frac{|\log \varepsilon|}{h_{\varepsilon}} \liminf_{\varepsilon \to 0} \frac{\pi \sum_i |d_i|}{h_{\varepsilon}} + \frac{1}{2} \liminf_{n \to \infty} \int_{\Omega \setminus \mathcal{A}_N} \frac{|\nabla h_{\varepsilon}|^2}{h_{\varepsilon}^2} + \left| \frac{h_{\varepsilon}}{h_{\varepsilon}} - 1 \right|^2.$$

Using (9.2), (9.7), the weak convergence of \( h_{\varepsilon}/h_{\varepsilon} \) to \( h_{\mu} \), and weak lower semicontinuity (up to a further extraction), we deduce that for every \( N \),

$$\liminf_{n \to \infty} \frac{G_{\varepsilon_n}(u_{\varepsilon_n}, A_{\varepsilon_n})}{h_{\varepsilon_n}^2} \geq \frac{1}{2\lambda} |\mu|(\Omega) + \frac{1}{2} \int_{\Omega \setminus \mathcal{A}_N} |\nabla h_{\mu}|^2 + |h_{\mu} - 1|^2.$$

Letting then \( N \rightarrow \infty \), since \( |\mathcal{A}_N| \rightarrow 0 \) we deduce that \( h_{\mu} \in L^2(\Omega) \) and the lower bound \( \liminf_{\varepsilon \to 0} \frac{G_{\varepsilon}(u_{\varepsilon}, A_{\varepsilon})}{h_{\varepsilon}} \geq E(\mu) \) holds along that subsequence. This proves the \( \Gamma \)-liminf relation (together with the compactness).
9.2.2 Upper bound

We prove the $\Gamma$-limsup inequality via the construction of a recovery sequence when $\mu$ is a nonnegative measure (the general case is not much different, see [SS4 Chap. 7] for details). We first split $G_\Omega(x,y)$ (defined by (7.16)) as $G_\Omega(x,y) = \frac{1}{\pi^2} (-\log|x-y| + S_\Omega(x,y))$ with $S_\Omega \in C^1(\Omega \times \Omega)$.

Step 1. Determining the vortex locations
Since $\mu$ is a positive measure of finite mass in $\Omega$, we may apply the $\Gamma$-limsup part of Proposition 2.8 to the probability measure $\mu \big|_{\mu |_\Omega}$, with potential $V = 0$ and with

$$n = \left\lfloor \frac{1}{2\pi} h_{ex} |\mu|_{\Omega} \right\rfloor,$$

(9.9)

with $\lfloor \cdot \rfloor$ the integer part. This yields the existence of points $a_i$ (depending on $n$ hence on $\varepsilon$), such that

$$\nu_\varepsilon := \sum_i \delta_{a_i} \to \mu \big|_{\mu |_\Omega} \quad (9.10)$$

and

$$\limsup_{n \to \infty} \iint_{\mathbb{R}^2 \setminus \Delta} -\log |x - y| d\nu_\varepsilon(x) d\nu_\varepsilon(y) \leq \frac{1}{|\mu|_{\Omega}} \iint -\log |x - y| d\mu(x) d\mu(y).$$

(9.11)

Moreover, examining the proof in Proposition 2.8, we see that the points $a_i$ are separated by a distance $Cn^{-1/2} \geq ch_{ex}^{-1/2} \gg \varepsilon$, we may also check that the points can be assumed to all lie in $\Omega$, and that the same results hold when replacing $\nu_\varepsilon$ by $\sum_i \mu_i h_{ex}$, where $\mu_i$ is the uniform measure of mass $2\pi$ supported in $\partial B(a_i, \varepsilon)$ (note the $\mu_i$'s have disjoint support by the previous observation). In other words, we have

$$\frac{\mu_\varepsilon}{h_{ex}} := \sum_i \frac{\mu_i}{h_{ex}} \to \mu$$

(9.12)

in the weak sense of measures, and

$$\limsup_{\varepsilon \to 0} \sum_{i \neq j} \frac{1}{h_{ex}} \iint -\log |x - y| d\mu_i(x) d\mu_j(y) \leq \iint -\log |x - y| d\mu(x) d\mu(y).$$

(9.13)

Since (9.12) holds, by weak convergence and regularity of $S_\Omega$ we also have

$$\limsup_{\varepsilon \to 0} \frac{1}{h_{ex}} \iint S_\Omega(x,y) d\mu_\varepsilon(x) d\mu_\varepsilon(y) \leq \iint S_\Omega(x,y) d\mu(x) d\mu(y).$$

(9.14)

We can also easily estimate the contributions of diagonal terms, by definition of $\mu_i$:

$$\sum_{i=1}^n \iint -\log |x - y| d\mu_i(x) d\mu_i(y) = -n \int_{[0,2\pi]^2} \log |\varepsilon e^{i\theta} - \varepsilon e^{i\phi}| d\theta d\phi$$

$$= 4\pi^2 n |\log \varepsilon| + C. \quad (9.15)$$
Combining (9.2), (9.9), (9.13)—(9.15) and the splitting of $G_\Omega$ we obtain

\[
\frac{1}{h_{\text{ex}}^2} \iint_{\Omega \times \Omega} G_\Omega(x, y) \, d\mu_\varepsilon(x) \, d\mu_\varepsilon(y) \leq \iint_{\Omega \times \Omega} G_\Omega(x, y) \, d\mu(y) + \frac{|\mu(\Omega)|}{\lambda} + o(1). \tag{9.16}
\]

**Step 2. Constructing the configuration.**

In this step we sort of reverse-engineer the configuration $(u_\varepsilon, A_\varepsilon)$ from the vortices we have constructed. First we let $h_\varepsilon$ be the solution of

\[
\begin{cases}
-\Delta h_\varepsilon + h_\varepsilon = \mu_\varepsilon & \text{in } \Omega \\
h_\varepsilon \geq h_{\text{ex}} & \text{on } \partial \Omega.
\end{cases} \tag{9.17}
\]

Then, we let $A_\varepsilon$ be any vector field such that $\text{curl} \, A_\varepsilon = h_\varepsilon$ in $\Omega$ and define $u_\varepsilon = \rho_\varepsilon e^{i\varphi_\varepsilon}$ as follows. We let

\[
\rho_\varepsilon(x) = \begin{cases}
0 & \text{in } \bigcup_{i=1}^n B(a_i, \varepsilon) \\
\frac{|x - a_i| - 1}{\varepsilon} & \text{in } B(a_i, 2\varepsilon) \setminus B(a_i, \varepsilon) \\
1 & \text{otherwise.}
\end{cases} \tag{9.18}
\]

For any $x \in \Omega \setminus \bigcup_{i=1}^n B(a_i, \varepsilon)$, we let

\[
\varphi_\varepsilon(x) = \oint_{(x_0, x)} (A_\varepsilon - \nabla^\perp h_\varepsilon) \cdot \tau \, d\ell,
\]

where $x_0$ is any reference point in $\Omega \setminus \bigcup_{i} B(a_i, \varepsilon)$, and $(x_0, x)$ is any curve joining $x_0$ to $x$ in $\Omega \setminus \bigcup_{i} B(a_i, \varepsilon)$. From (9.17) we see that this definition does not depend modulo $2\pi$ on the curve chosen to join $x_0$ to $x$, thus $e^{i\varphi_\varepsilon}$ is well-defined in $(\cup_i B(a_i, \varepsilon))^c$. Indeed, if $\gamma = \partial U$ is a closed curve in $(\cup_i B(a_i, \varepsilon))^c$, using Stokes’ theorem and $\text{curl} \, A_\varepsilon = h_\varepsilon$, we find

\[
\oint_\gamma (A_\varepsilon - \nabla^\perp h_\varepsilon) \cdot \tau \, d\ell \overset{\text{Stokes}}{=} \int_U (-\Delta h_\varepsilon + h_\varepsilon) = \int_U \sum_i \mu_i \in 2\pi \mathbb{N}.
\]

The (multi-valued) function $\varphi_\varepsilon$ satisfies

\[
-\nabla^\perp h_\varepsilon = \nabla \varphi_\varepsilon - A_\varepsilon \quad \text{in } \Omega \setminus \bigcup_{i} B(a_i, \varepsilon). \tag{9.19}
\]

Finally, we may define $u_\varepsilon = \rho_\varepsilon e^{i\varphi_\varepsilon}$ and we notice that the fact that $\varphi_\varepsilon$ is not defined on $\cup_i B(a_i, \varepsilon)$ is not important since $\rho_\varepsilon$ is zero there.

**Step 3. Computing the energy of the test-configuration.**

To finish it suffices to evaluate $G_\varepsilon(u_\varepsilon, A_\varepsilon)$. First we notice that by construction

\[
\int_{\Omega} \nabla |u_\varepsilon|^2 + \frac{(1 - |u_\varepsilon|^2)^2}{2\varepsilon^2} \leq Cn \leq o(h_{\text{ex}}^2)
\]
The leading order behavior for Ginzburg-Landau

by (9.9). Using that \(|\nabla A_\varepsilon u_\varepsilon|^2 = |\nabla u_\varepsilon|^2 + |u_\varepsilon|^2 |\nabla \varphi_\varepsilon - A_\varepsilon|^2\) and that \(|u_\varepsilon| = 0\) in \(\cup_i B(a_i, \varepsilon)\) and \(|u_\varepsilon| \leq 1\) everywhere and (9.19), we deduce that

\[
\int_\Omega |\nabla A_\varepsilon u_\varepsilon|^2 + \frac{(1 - |u_\varepsilon|^2)^2}{2 \varepsilon^2} \leq \int_{(\cup_i B(a_i, \varepsilon))} |\nabla \varphi_\varepsilon - A_\varepsilon|^2 + O(\varepsilon) \leq \int_\Omega |\nabla h_\varepsilon|^2 + o(h_{ex}^2).
\]

It follows that for this configuration we have the inequality corresponding to (1.2) i.e.

\[
G_\varepsilon(u_\varepsilon, A_\varepsilon) \leq \frac{1}{2} \int_\Omega |\nabla h_\varepsilon|^2 + |h_\varepsilon - h_{ex}|^2 + o(h_{ex}^2).
\]

In view of (9.17), we have \(h_\varepsilon(x) = h_{ex} + \int_\Omega G_\Omega(x, y)(\mu_\varepsilon(y) - h_{ex})\, dy\) and using (9.1), we have that

\[
\int_\Omega |\nabla h_\varepsilon|^2 + |h_\varepsilon - h_{ex}|^2 = \int_{\Omega \times \Omega} G_\Omega(x, y)\, d(\mu_\varepsilon(x) - h_{ex})(x)\, d(\mu_\varepsilon(y) - h_{ex})(y),
\]

as in (1.4). Evaluating this integral is now a direct consequence of (9.16) and (9.12) and leads us to

\[
\limsup_{\varepsilon \to 0} \frac{G_\varepsilon(u_\varepsilon, A_\varepsilon)}{h_{ex}^2} \leq \frac{1}{2} \int_{\Omega \times \Omega} G_\Omega(x, y)\, d(\mu - 1)(x)\, d(\mu(y) - 1)(y) + \frac{1}{2\lambda} |\mu|(\Omega)
\]

\[
= \frac{1}{2} \int_\Omega |\nabla h_{\mu}|^2 + |h_{\mu} - 1|^2 + \frac{1}{2\lambda} |\mu|(\Omega) = E_\lambda(\mu).
\]

Indeed, there is no problem in passing to the limit in terms of the form

\[
\int_{\Omega \times \Omega} G_\Omega(x, y)\, dx\, d\frac{h_\varepsilon}{h_{ex}}(y)
\]

since one may check that the function \(\int_\Omega G_\Omega(x, y)\, dx\) is a continuous function of \(y\). This concludes the proof of the \(\Gamma\)-limsup, provided we check that we do have \(\mu(u_\varepsilon, A_\varepsilon) \Rightarrow \mu\). But this can be checked from \(\mu(u_\varepsilon, A_\varepsilon) = \text{curl} (|u_\varepsilon|^2(\nabla \varphi_\varepsilon - A_\varepsilon)) + h_\varepsilon = \text{curl} (|u_\varepsilon|^2 \nabla h_\varepsilon) + h_\varepsilon, \quad (9.17), (9.18)\) and (9.12).

9.3 Minimization of the mean-field limit and connection to the obstacle problem

Once the \(\Gamma\)-convergence result is obtained, it immediately implies the leading-order behavior from Proposition 2.6 and Remark 2.7, we have

**Corollary 9.2** (Limit of Ginzburg-Landau minimizers). Assume (9.2). Let \((u_\varepsilon, A_\varepsilon)\) minimize \(G_\varepsilon\), then as \(\varepsilon \to 0\) we have

\[
\frac{\mu(u_\varepsilon, A_\varepsilon)}{h_{ex}} \Rightarrow \mu_\lambda
\]

where \(\mu_\lambda\) is the unique minimizer of \(E_\lambda\).
The fact that $E_\lambda$ has a unique minimizer is a consequence of its obvious convexity in $\mu$. This result is indeed a mean-field limit, since it describes the limit of the (suitably normalized) vorticity for which (7.12) holds. Since $h_{ex} \to +\infty$ as $\varepsilon \to 0$, the number of vortices is expected to blow-up like $h_{ex}$ too, and they arrange themselves according to the distribution $\mu_\lambda$, which plays the role of the equilibrium measure $\mu_0$ for the Coulomb gas in the Ginzburg-Landau context, cf. Figure 9.1.

The limiting energy $E_\lambda(\mu)$ is of similar nature as the mean-field limit Hamiltonian $I$ in Chapter 2. This is more readily visible if one rewrites $E_\lambda$ as

$$E_\lambda(\mu) = \frac{1}{2\lambda} |\mu|(\Omega) + \int_{\Omega \times \Omega} G_\Omega(x, y)d(\mu - 1)(x) d(\mu - 1)(y).$$

(9.20)

Compared to $I$ in (2.7), the confining potential is replaced by the fact of working in a bounded domain with a Green-Dirichlet function, and the constraint that $\mu$ be a probability is replaced by the penalization term in $|\mu|(\Omega)$, which behaves like a Lagrange multiplier term.

We may now identify the minimizer of $E_\lambda$ with the solution of an obstacle problem, just like for the minimization of $I$. The correspondence here is in some sense even easier due to the fact that we are in a bounded domain.

**Proposition 9.3** (Identification of the optimal density). The minimizer $\mu_\lambda$ of $E_\lambda$ is uniquely characterized by the fact that the associated potential $h_{\mu_\lambda}$ given by (9.4) is the solution of the following obstacle problem :

$$\min_{\begin{array}{c} h \geq 1 - \frac{1}{2\lambda} \\ h - 1 \in H_0^1(\Omega) \end{array}} \int_{\Omega} |\nabla h|^2 + h^2.$$  

(9.21)

The function $h_{\mu_\lambda}$ is in turn characterized by the variational inequality

$$\forall v \geq 1 - \frac{1}{2\lambda}, v - 1 \in H_0^1(\Omega), \text{ we have } \int_{\Omega} \nabla h_{\mu_\lambda} \cdot \nabla (v - h_{\mu_\lambda}) + h_{\mu_\lambda} (v - h_{\mu_\lambda}) \geq 0$$

or by the relations

$$\begin{cases} h_{\mu_\lambda} \geq 1 - \frac{1}{2\lambda} & \text{in } \Omega \\ h_{\mu_\lambda} = 1 - \frac{1}{2\lambda} & \text{q.e. in the support of } \mu_\lambda \\ h_{\mu_\lambda} = 1 & \text{on } \partial \Omega. \end{cases}$$

(9.22)

The constant function $1 - \frac{1}{2\lambda}$ thus plays the role of the obstacle. We will write

$$\omega_\lambda := \left\{ h_{\mu_\lambda} = 1 - \frac{1}{2\lambda} \right\}$$

(9.23)

for the coincidence set. From the above characterizations we deduce that $\mu_\lambda$ is a nonnegative measure, and that

$$\mu_\lambda = -\Delta h_{\mu_\lambda} + h_{\mu_\lambda} = (1 - \frac{1}{2\lambda})1_{\omega_\lambda}.$$  

(9.24)
Thus, in this Ginzburg-Landau context, the optimal measure always has a density, and that density is always constant on its support, cf. Fig. 9.1.

Figure 9.1. The optimal measure for Ginzburg-Landau

**Proof of the proposition.** To give an alternate proof to that of Chapter 2 we may obtain this by convex duality. It suffices to observe that the minimization of $E_\lambda$ viewed as the following function of $h$:

$$
\mu = \frac{1}{2\lambda} \int_{\Omega} |\nabla h|^2 + h^2
$$

is dual in the sense of convex duality to the minimization problem

$$
\min_{\{h \leq \frac{1}{2}\lambda\}} \frac{1}{2} \int_{\Omega} |\nabla h|^2 + h^2 + 2h
$$

in the sense that they have the same minimizer (and minimum). For a proof of this fact, cf. [SS4, Chap. 7]. One can then check that by the maximum principle the solution $h$ satisfies $h \leq 0$, so the upper constraint $h \leq \frac{1}{2}\lambda$ is not active, and thus $h + 1$ solves (9.21). The alternate way is to start from (9.20) and make variations on $\mu$ as in the proof of Theorem 2.1. This leads to the equations (9.22), and one can then check that they uniquely characterize the solution to (9.21).

The regularity theory is exactly as in Proposition 2.22: since the obstacle is constant hence smooth, $h_{\mu_\lambda}$ is $C^{1,1}$ by Frehse’s regularity theorem, and we can deduce that (9.24) holds.

The solution of (9.21) when the constraint of being above the obstacle is omitted, is obviously the function $h_0$ solution to (7.18). It then follows that $h_0$ is also the solution of the problem (9.21), if and only if $h_0$ lies above the obstacle i.e. $h_0 \geq 1 - \frac{1}{2\lambda}$, equivalent to $\lambda \geq (2 \min(h_0 - 1))$. Whether this condition is
9.4 The intermediate regime near $H_{c_1}$

satisfied depends on the value of $\lambda$, which we recall is $\lim_{\varepsilon \to 0} \frac{h_{\text{ex}}}{|\log \varepsilon|}$, i.e. encodes the intensity of the applied magnetic field.

We deduce the following result on the description of $\mu_\lambda$ as the external magnetic field is increased.

**Proposition 9.4.**

- $\omega_\lambda$ is increasing with respect to $\lambda$ and $\cup_{\lambda > 0} \omega_\lambda = \Omega$.
- For $\lambda \leq \lambda_\Omega := \frac{1}{2 \max |h_0 - 1|}$ we have $\omega_\lambda = \emptyset$, $\mu_\lambda = 0$ and $h_{\mu_\lambda} = h_0$.
- For $\lambda > \lambda_\Omega := \frac{1}{2 \max |h_0 - 1|}$ we have $\mu_\lambda \neq 0$, and (9.24) holds.

This way we recover in a weaker sense the value of the first critical field for which vorticity first appears:

$$H_{c_1} \sim \lambda_\Omega |\log \varepsilon|$$  \hspace{1cm} (9.25)

i.e. we give a first rigorous justification of (7.20).

We can also deduce from the proposition that the onset of $\omega_\lambda$ is located near the point(s) of minimum of $h_0$, i.e. that the vorticity first appears there, as formally derived in Chapter 7.

9.4 The intermediate regime near $H_{c_1}$

Understanding what happens more precisely near $H_{c_1}$ (exact number and locations of the vortices as $h_{\text{ex}} \sim H_{c_1}$) requires a finer analysis than this leading order one: one needs to make more precise expansions around $h_{\text{ex}}h_0$ as in Section 7.4.2. This is done in [SS4, Chap. 9, Chap. 12]. It is found that the vortices appear one by one near the point(s) of minimum of $h_0$ in $\Omega$, with new vortices appearing each time $h_{\text{ex}}$ is incremented by an order $\log |\log \varepsilon|$, as long as $h_{\text{ex}} \sim \frac{|\log \varepsilon|}{2 \max |h_0 - 1|}$ i.e. $\lambda = 1/(2 \max |h_0 - 1|)$. The locations of the vortices tend to minimize exactly a Coulomb gas type of interaction: if their number $n$ remains bounded as $\varepsilon \to 0$, their location (suitably blown-up near the point(s) of minimum of $h_0$) minimize

$$- \sum_{i \neq j \in [1,n]} \log |x_i - x_j| + n \sum_{i=1}^{n} Q(x_i)$$  \hspace{1cm} (9.26)

while if their number becomes unbounded, their density (again after suitable scaling) tends to minimize among probability measures

$$I(\mu) = \int - \log |x - y| d\mu(x) d\mu(y) + \int_{\mathbb{R}^2} Q(x) d\mu(x)$$  \hspace{1cm} (9.27)

with $Q$ a nonnegative quadratic function equal to the Hessian of $h_0$ at the point of minimum of $h_0$. Once $\lambda > 1/(2 \max |h_0 - 1|)$ the optimal description is the one just given in Propositions 9.3, 9.4.
In other words, in the critical regime near $H_{c_1}$, the interaction of vortices is precisely that of a 2D Coulomb gas with quadratic confining potential $Q$. One remembers from Chapter 2 that the equilibrium measures associated to quadratic potentials $V$ are always of constant density, just like what happens for $\mu_\lambda$. 
10 The splitting and the next order behavior for Ginzburg-Landau

In this chapter we sketch the method that allows us to derive the renormalized energy $W$ from the minimization of Ginzburg-Landau, at the next order, beyond the mean-field limit that we just saw in Chapter 9. This is, in a simplified form, the content of the paper [SS6].

10.1 Splitting

In order to extract the next order energy, it is very important to have an exact splitting of the Ginzburg-Landau functional which “algebraically” decouples the orders. We saw in the previous chapter that the magnetic field $h_\varepsilon$ satisfies $h_\varepsilon \rightarrow h_{\mu_\lambda}$ in some weak sense, with $\mu_\lambda$ the “mean-field” limit, i.e. the minimizer of $E_\lambda$. Since $h_\varepsilon$ plays the role of the potential $h_n$ for the Coulomb gas, the splitting for the Coulomb gas Hamiltonian in Chapter 3 can then give us a hint: to split $h_\varepsilon$ as $h_{\text{ex}} h_{\mu_\lambda} + h_1$ where $h_1$ is a remainder. This rough idea is correct, however there are two difficulties: first (1.2) is only an approximate relation, and not an identity, so we really need to work starting from the configurations $(u, A)$ themselves. Secondly the approximation $h \sim h_{\text{ex}} h_{\mu_\lambda}$ is correct at leading order, but a correction needs to be introduced in order to extract the right next order.

The fact is that we need to know the number of vortices (or their total degree) more precisely than through $h_{\text{ex}} h_{\mu_\lambda}$. $\mu_\lambda$ was found by minimizing the limiting energy $E_\lambda$, which was itself derived by bounding below the cost of the self-interaction of each vortex via the ball construction lower bound, which gives a cost $\sim \pi |d||\log \varepsilon|$ per vortex, resulting after taking the limit $\varepsilon \rightarrow 0$ in the term $\lambda |\mu|(\Omega)$ in $E_\lambda$. But one needs to be more precise: we will restrict ourselves to the situation where $\lambda > \lambda_\Omega$ (cf. (9.25)), i.e. the limiting measure $\mu_\lambda$ and the coincidence set $\omega_\lambda$ are nontrivial. The number of expected vortices in that regime is thus proportional to $h_{\text{ex}}$. Since vortices are uniformly distributed in $\omega_\lambda$, their mutual distances are of order $1/\sqrt{h_{\text{ex}}}$. We can thus think of each vortex as being alone in a box of size $C/\sqrt{h_{\text{ex}}}$, and remembering that its core size is always $\varepsilon$, a lower bound of the type $\frac{\lambda}{\Omega}$ leads us to expecting a cost $\pi |d||\log C\sqrt{h_{\text{ex}}}$ per vortex. In other words, there should be a correction of order $\log h_{\text{ex}} \sim \log |\log \varepsilon|$ per vortex, which has not been accounted for in $E_\lambda$ (and did not matter at the leading order level). This heuristically justifies introducing this correction in the self-interaction cost, by minimizing instead of $E_\lambda$ the following problem:

$$
\frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \int_\Omega |\mu| + \frac{1}{2} \int_\Omega |\nabla h|^2 + |h - h_{\text{ex}}|^2
$$

(10.1)
where $\mu$ and $h$ are related via

$$
\begin{cases}
-\Delta h + h = \mu & \text{in } \Omega \\
h = h_{\text{ex}} & \text{on } \partial \Omega.
\end{cases}
$$

(10.2)

This is a problem of the same form as the minimization of $E_\lambda$ (to see it, just divide everything by $h_{\text{ex}}^2$), except that the parameter $\lambda$ is replaced by the correction \( \frac{h_{\text{ex}}}{\log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}}}, \) which is equivalent to $\lambda$ as $\varepsilon \to 0$ (but here we define this energy for each fixed $\varepsilon$). As a result, the same proof as Proposition 9.3 applies, and asserts that the minimization problem (10.1) is equivalent to the obstacle problem

$$
\min_{h \geq h_{\text{ex}} - \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}}} \int_\Omega |\nabla h|^2 + h^2.
$$

(10.3)

We will denote the solution by $h_{0,\varepsilon}$ and the associated measure

$$
\mu_{0,\varepsilon} = -\Delta h_{0,\varepsilon} + h_{0,\varepsilon}.
$$

(10.4)

It is clear that $\frac{h_{\text{ex}}}{\lambda} \to \mu_\lambda$ and $\frac{h_{\text{ex}}}{h_{0,\varepsilon}} \to h_{\mu_\lambda}$ as $\varepsilon \to 0$, however these objects are a little more precise, and as explained above, they are the ones with respect to which we should do the splitting.

We will also denote by $\omega_{0,\varepsilon} = \{x, h_{0,\varepsilon}(x) = h_{\text{ex}} - \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}}\}$ the corresponding coincidence set, and note that

$$
\mu_{0,\varepsilon} = \left( h_{\text{ex}} - \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \right) 1_{\omega_{0,\varepsilon}},
$$

and recall that $\mu_{0,\varepsilon} \geq 0$.

**Proposition 10.1** (Splitting formula for Ginzburg-Landau [SS6]). Let $(u, A)$ be an arbitrary configuration and for any $\varepsilon$ and $h_{\text{ex}}$, set

$$
A_{1,\varepsilon} = A - \nabla \perp h_{0,\varepsilon},
$$

where $h_{0,\varepsilon}$ is the solution of (10.3). Then we have

$$
G_\varepsilon(u, A) = G_0^{\varepsilon} + G_1^{\varepsilon}(u, A_{1,\varepsilon}) - \frac{1}{2} \int_\Omega (1 - |u|^2)|\nabla h_{0,\varepsilon}|^2
$$

(10.5)

where

$$
G_0^{\varepsilon} = \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \int_\Omega \mu_{0,\varepsilon} + \frac{1}{2} \int_\Omega |\nabla h_{0,\varepsilon}|^2 + |h_{0,\varepsilon} - h_{\text{ex}}|^2
$$

(10.6)

and

$$
G_1^{\varepsilon}(u, A) = \frac{1}{2} \int_\Omega |\nabla A|^2 + |\text{curl } A - \mu_{0,\varepsilon}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2} + \int_\Omega (h_{0,\varepsilon} - h_{\text{ex}}) \mu(u, A).
$$
Proof of the splitting formula. First, in view of the definition of $A_{1,\varepsilon}$ and (10.4), we may write

$$|
abla A| = |\nabla A_{1,\varepsilon}| + |\nabla h_{0,\varepsilon}| - 2 \nabla \perp h_{0,\varepsilon} \cdot \langle i u, \nabla A_{1,\varepsilon} u \rangle,$$

and

$$\text{curl} A = \text{curl} A_{1,\varepsilon} + \Delta h_{0,\varepsilon} = \text{curl} A_{1,\varepsilon} + h_{0,\varepsilon} - \mu_{0,\varepsilon}.$$

Inserting into the expression of $G_{\varepsilon}$ and expanding the squares, we find

$$G_{\varepsilon}(u, A) = \frac{1}{2} \int_{\Omega} |u|^2 |\nabla h_{0,\varepsilon}|^2 + |h_{0,\varepsilon} - h_{\text{ex}}|^2$$

$$+ \int_{\Omega} -\nabla \perp h_{0,\varepsilon} \cdot \langle i u, \nabla A_{1,\varepsilon} u \rangle + (\text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon})(h_{0,\varepsilon} - h_{\text{ex}})$$

$$+ \frac{1}{2} \int_{\Omega} |\nabla A_{1,\varepsilon} u|^2 + |\text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2}$$

$$= \frac{1}{2} \int_{\Omega} |\nabla h_{0,\varepsilon}|^2 + |h_{0,\varepsilon} - h_{\text{ex}}|^2 + \int_{\Omega} (h_{0,\varepsilon} - h_{\text{ex}})(\text{curl} \langle i u, \nabla A_{1,\varepsilon} u \rangle + \text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon})$$

$$+ \frac{1}{2} \int_{\Omega} (|u|^2 - 1)|\nabla h_{0,\varepsilon}|^2$$

where we have used an integration by parts and the fact that $h_{0,\varepsilon} = h_{\text{ex}}$ on $\partial \Omega$.

We next observe that

$$\int_{\Omega} (h_{0,\varepsilon} - h_{\text{ex}}) \mu_{0,\varepsilon} = -\frac{1}{2} \log \frac{1}{\varepsilon h_{\text{ex}}} \int_{\Omega} \mu_{0,\varepsilon}$$

since that is the value of $h_{0,\varepsilon} - h_{\text{ex}}$ on the support of $\mu_{0,\varepsilon}$, and

$$\text{curl} \langle i u, \nabla A_{1,\varepsilon} u \rangle + \text{curl} A_{1,\varepsilon} = \mu(u, A_{1,\varepsilon}).$$

Inserting into the above, we obtain the result. □

As desired, we have obtained an exact splitting formula for the Ginzburg-Landau energy. The first term $G_{\varepsilon}^0$ is a constant independent of $(u, A)$ and easily seen to be asymptotically equivalent to $h_{\text{ex}}^2 E_{\lambda}(\mu_{\lambda})$, i.e. to the leading order of the energy. The last term is generally $o(1)$ because thanks to the potential term in the energy we may control $\int_{\Omega}(1 - |u|^2)$ by $\varepsilon \sqrt{G_{\varepsilon}(u, A)}$ via Cauchy-Schwarz. The middle term $G_{\varepsilon}^1$ is the interesting one: it is the difference between an energy functional which is very similar to Ginzburg-Landau, except with external field replaced by the non constant function $\mu_{0,\varepsilon}$, and a term which, thanks to the Jacobian estimate (Theorem 8.2) can be evaluated by

$$\int_{\Omega} (h_{0,\varepsilon} - h_{\text{ex}}) \mu(u, A) \simeq 2\pi \sum_i d_i (h_{0,\varepsilon} - h_{\text{ex}})(a_i)$$

(indeed, one may easily check that $\mu(u, A_{1,\varepsilon}) \simeq \mu(u, A)$). In this term, all vortices with positive degree bring a negative contribution, since $h_{0,\varepsilon} \leq h_{\text{ex}}$ in $\Omega$ by the
The splitting and the next order behavior for Ginzburg-Landau maximum principle. In other words they will allow to gain energy, while negative degree vortices will not and hence will not be favorable. Moreover, \( h_{0,\varepsilon} - h_{\text{ex}} \) is minimal and equal to \(-\frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}}\) in the coincidence set \( \omega_{0,\varepsilon} \), hence it will be most favorable to have vortices there. Setting

\[
\zeta_{\varepsilon}(x) = h_{0,\varepsilon} - h_{\text{ex}} + \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}}
\]

(10.7)

we have \( \zeta_{\varepsilon} \geq 0 \) in \( \Omega \), \( \{ \zeta_{\varepsilon} = 0 \} = \omega_{0,\varepsilon} \), and we may thus rewrite the splitting formula formally as

\[
G_{\varepsilon}(u, A) \simeq G_{\varepsilon}^0 + 2\pi \sum_i d_i \zeta_{\varepsilon}(a_i)
\]

\[
+ \frac{1}{2} \int_{\Omega} |\nabla A_{1,\varepsilon}u|^2 + |\text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2} - \pi n \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} + o(1),
\]

(10.8)

with \( n = \sum_i d_i \) is the number of vortices (assumed positive), and rigorously as

\[
G_{\varepsilon}(u, A) = G_{\varepsilon}^0 + \int_{\Omega} \zeta_{\varepsilon} \mu(u, A_{1,\varepsilon}) + F_{\varepsilon}(u, A_{1,\varepsilon}) + o(1)
\]

(10.9)

where

\[
F_{\varepsilon}(u, A_{1,\varepsilon}) := \frac{1}{2} \int_{\Omega} |\nabla A_{1,\varepsilon} u|^2 + |\text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2}
\]

\[- \frac{1}{2} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \int_{\Omega} \mu(u, A_{1,\varepsilon}).
\]

(10.10)

This is of the same form as the splitting of the Coulomb gas Hamiltonian \( H_n \) in Proposition 3.4. The role of \( n^2 I(\mu_0) \) is played by \( G_{\varepsilon}^0 \sim h_{\text{ex}}^{-2} E_\lambda(\mu_\lambda) \), \( \zeta \) plays the same role of a confining potential as \( \zeta \), confining the points to the support of the optimal measure \( (\mu_{0,\varepsilon} \text{ for Ginzburg-Landau, the equilibrium measure for the Coulomb gas}) \), and plays no role otherwise. The remaining term \( F_{\varepsilon} \) behaves as the precursor to the renormalized energy \( W(\nabla h_n, 1_{\mathbb{R}^2}) \), although this is more delicate to see: the term \( \frac{1}{2} \int_{\Omega} |\nabla A_{1,\varepsilon} u|^2 + |\text{curl} A_{1,\varepsilon} - \mu_{0,\varepsilon}|^2 + \frac{(1 - |u|^2)^2}{2\varepsilon^2} \) behaves like a Ginzburg-Landau energy, hence it will include the interaction between vortices, plus the cost of each vortex, which can be estimated via a ball-construction method, while the term \( -\pi n \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \) in effect subtracts off the cost of each vortex, i.e. “renormalizes” the Ginzburg-Landau energy.

All of the analysis from that point consists in showing rigorously that this is true, and that the term \( F_{\varepsilon} \) will effectively converge to the (average of the) renormalized energy \( W \). This will be more technical than for the Coulomb gas, because for instance we have to truly get rid of the possibility of (too many) negative vortices. Also we will need to have very precise ball construction lower bounds for the cost of each vortex, to show that the compensation \( -\pi n \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}} \), which includes the correction in \log \log \varepsilon \) that we inserted, is the right one.
Remark 10.2. In [SS6] the computations are made more complicated by the fact that we also treat the case of $h_{\text{ex}}$ possibly very close to $H_{c1}$, which requires more precise estimates, themselves requiring the mass of the measure which respect to which one splits to be quantized.

10.2 Deriving $W$ from Ginzburg-Landau

10.2.1 Rescaling and notation

In view of the splitting formula (10.5), and the fact that the last term is very small, in order to study $G_{\varepsilon}$, it suffices to study $G_{1,\varepsilon}^\varepsilon(u,A_{1,\varepsilon})$. We may introduce $h_{1,\varepsilon} = \text{curl} A - h_{0,\varepsilon}$. In view of the fact (which we may assume) that for $(u,A)$ the second Ginzburg-Landau equation (7.13), hence the London equation (7.14) is satisfied, we check that (by definition of $A_{1,\varepsilon}$ and $h_{0,\varepsilon}$), the function $h_{1,\varepsilon}$ satisfies

\[
\begin{cases}
-\Delta h_{1,\varepsilon} + h_{1,\varepsilon} = \mu(u,A) - \mu_{0,\varepsilon} & \text{in } \Omega \\
\quad h_{1,\varepsilon} = 0 & \text{on } \partial \Omega.
\end{cases}
\] (10.11)

Thus $h_{1,\varepsilon}$ is the analogue in the Ginzburg-Landau context of the potential $h_n$ defined in (3.5).

As in the case of the Coulomb gas, the next step is to blow up at the scale of the inter-vortex distance, here of order $1/\sqrt{h_{\text{ex}}}$. Figure 10.1 illustrates how we blow up around a center point belonging to the coincidence set of the obstacle problem, i.e. the support of $\mu_{\lambda}$ (or $\mu_{0,\varepsilon}$), and wish to find a triangular lattice distribution of vortices after blow-up in the limit $\varepsilon \to 0$.

We thus define

\[x' = \sqrt{h_{\text{ex}}} x, \quad \varepsilon' = \sqrt{h_{\text{ex}}} \varepsilon, \quad u'(x') = u(x),\]

\[A'(x') = \sqrt{h_{\text{ex}}} A_{1,\varepsilon}(\frac{x'}{\sqrt{h_{\text{ex}}}}), \quad h'(x') = h_{1,\varepsilon}(\frac{x'}{\sqrt{h_{\text{ex}}}}), \quad \mu'(x') = \mu(u,A)(\frac{x'}{\sqrt{h_{\text{ex}}}})\]

\[\Omega' = \sqrt{h_{\text{ex}}} \Omega, \quad \omega_{0,\varepsilon} = \sqrt{h_{\text{ex}}} \omega_{0,\varepsilon}, \quad \mu_{0,\varepsilon} = (1 - \frac{1}{2h_{\text{ex}}} \log \frac{1}{\varepsilon \sqrt{h_{\text{ex}}}})1_{\omega_{0,\varepsilon}}.\]

We note that the density of $\mu'_{0,\varepsilon}$ tends to $m_{\lambda} := 1 - \frac{1}{2\lambda}$ as $\varepsilon \to 0$. Rescaling the equation (10.11) yields

\[-\Delta h' + \frac{1}{h_{\text{ex}}} h' = \mu' - \mu_{0,\varepsilon}\]

and thus if one centers the blow-up in a point of $\omega_{0,\varepsilon}$ we will have in the limit $\varepsilon \to 0$

\[-\Delta h = 2\pi \sum_p \delta_p - m_{\lambda} \text{ in } \mathbb{R}^2\]

for some points $p \in \mathbb{R}^2$ (which we will prove appear with single-multiplicity), i.e, $\nabla h$ belongs to the admissible class $\mathcal{A}_{m_{\lambda}/2\pi}$ defined in Chapter 5, for which we can define $W(\nabla h)$. 

Inserting this change of scales into the energy $F_{\varepsilon}$, we find

$$F_{\varepsilon}(u, A_{1,\varepsilon}) = \frac{1}{2} \int_{\Omega'} \left| \nabla A' u' \right|^2 + \frac{1}{2} \left| \nabla A' \right|^2 + \frac{1}{2} \left| \nabla A' - \mu'_{0,\varepsilon} \right|^2 \right| + \frac{1 - |u'|^2}{2(e')^2} - \frac{1}{2} \log \left( \frac{1}{e'} \right) + \frac{1}{2} \int_{\Omega'} \mu(u', A').$$

As always, the main result will be obtained by proving first a lower bound, and second a matching upper bound via an explicit construction.

**10.2.2 Lower bound**

It is to be obtained by the abstract method presented in Section 6.1 applied on the large sets $\omega'_{0,\varepsilon}$, with the “local” energy being naturally

$$f_{\varepsilon}(u, A) = \int_{\mathbb{R}^2} \chi \left[ \frac{1}{2} \left| \nabla A u \right|^2 + \frac{h_{\text{ex}}}{2} \left| \nabla A - \mu_{0,\varepsilon} \right|^2 + \frac{1 - |u|^2}{4(e')^2} - \frac{1}{2} \log \left( \frac{1}{e'} \right) \mu(u, A) \right],$$

with $\chi$ some smooth nonnegative cutoff function supported in $B(0, 1)$ and of integral 1. Note that this local energy does not depend on a centering point, it is translation-invariant, so we may apply the method of Section 6.1 dropping the dependence in the centering point $x$. The energy outside of the coincidence set $\omega_{0,\varepsilon}$ will simply be discarded, it is indeed negligible for minimizers.
10.2 Deriving $W$ from Ginzburg-Landau

In order to apply the abstract framework of Section 6.1, a first step is to show that if
\[ \forall R > 0, \int_{K_R} f_\varepsilon(\theta_\lambda(u_\varepsilon,A_\varepsilon)) \, d\lambda \leq C_R \quad (10.13) \]
then $(u_\varepsilon, A_\varepsilon)$ has a subsequence converging to some $(u, A)$ with
\[ \liminf_{\varepsilon \to 0} f_\varepsilon(u_\varepsilon, A_\varepsilon) \geq f(u, A). \]

We also recall that (10.13) is equivalent to
\[
\int_{\mathbb{R}^2} \chi_{K_R} \left[ \frac{1}{2} |\nabla A u|^2 + \frac{h_{\text{ex}}}{2} |\text{curl} A - \mu_0'|^2 + \frac{(1 - |u|^2)^2}{4(\varepsilon')^2} - \frac{1}{2} \log \varepsilon' |\mu(u, A)| \right] \leq C_R.
\]

One of the important steps of the proof is to show that the energy density $f_\varepsilon$ controls the number of vortices, so that an upper bound of the form (10.13) implies that the total degree of vortices in $K_{R-1}$ is bounded by a constant (depending on $C_R$). Such a bound then easily implies that
\[
\int_{K_{R-1}} \frac{1}{2} |\nabla A u|^2 + \frac{h_{\text{ex}}}{2} |\text{curl} A - \mu_0'|^2 + \frac{(1 - |u|^2)^2}{4(\varepsilon')^2} \leq C_R |\log \varepsilon'|.
\]

As explained in Chapter 8, such an upper bound, which controls the number of vortices independently of $\varepsilon'$ makes our life much easier, since it puts us in the framework of [BBH, BR2], for which we can compute sharp and precise lower bounds (up to $o(1)$) for the Ginzburg-Landau energy. It also completely rules out the possibility of vortices with degrees different from +1. This type of analysis leads to the following lower bound: if (10.13) holds with $C_R$ replaced by $C_R^2$ and the second Ginzburg-Landau equation (7.13) holds, then up to extraction of a subsequence, we have that $h'$ converges to some $h \in A_{m_{\lambda}/2\pi}$ and
\[
\liminf_{\varepsilon \to 0} f_\varepsilon(u_\varepsilon, A_\varepsilon) \geq W(\nabla h, \chi) + \frac{\gamma}{2\pi} m_\lambda,
\]
where $\gamma$ is the constant from [BBH, Mi], and $W$ is the precursor to the renormalized energy as in Definition 3.1. The heuristic for this is quite natural: once the vortices, at points $p$, have been shown to be of degree 1 and in bounded number, we split the positive part of the energy as
\[
\int_{\mathbb{R}^2} \chi \left[ \frac{1}{2} |\nabla A u|^2 + \frac{h_{\text{ex}}}{2} |\text{curl} A - \mu_0'|^2 + \frac{(1 - |u|^2)^2}{4(\varepsilon')^2} \right]
\geq \int_{\mathbb{R}^2 \setminus \bigcup B(p,\varepsilon M)} \frac{1}{2} \chi |\nabla A u|^2 + \int_{\bigcup B(p,\varepsilon M) \setminus B(p,\varepsilon M')} \frac{1}{2} \chi |\nabla A u|^2
\]
\[ + \int_{\bigcup B(p,\varepsilon M')} \chi \left[ \frac{1}{2} |\nabla A u|^2 + \frac{h_{\text{ex}}}{2} |\text{curl} A - \mu_0'|^2 + \frac{(1 - |u|^2)^2}{4(\varepsilon')^2} \right]. \]

Outside the $B(p,\varepsilon M)$ with $M$ large, we expect that $|u| \simeq 1$ (which allowed us to discard some positive terms expected to be negligible). Then from the second
The splitting and the next order behavior for Ginzburg-Landau

Ginzburg-Landau equation, as seen in Chapter 7, we have $|\nabla \Delta u| \approx |\nabla h|$. In the annuli $B(p, \eta) \setminus B(p, M_\varepsilon)$, the energy is then expected to be bounded from below as in (8.3), which yields $\pi \log \frac{\eta}{\varepsilon}$ per vortex. In the “vortex cores” $B(p, M_\varepsilon)$ all the energy terms will matter, and the energy depends on the optimal radial profile for the Ginzburg-Landau energy, as given in [HH, Mi], which gives a contribution $\gamma$ per vortex, in addition to the cost $\pi \log M$. Combining all these terms, and multiplying by the number of vortices, which is expected to be $m_\lambda/2\pi$ per unit volume, we formally get

$$f_\varepsilon(u, A) \geq \int_{\mathbb{R}^2 \setminus \cup_p B(p, \eta)} \chi|\nabla h|^2 + \frac{\gamma}{2\pi} m_\lambda + \sum_p \chi(p) \left( \pi \log \frac{\eta}{\varepsilon} - \pi \log \frac{1}{\varepsilon} \right)$$

$$= \int_{\mathbb{R}^2 \setminus \cup_p B(p, \eta)} \chi|\nabla h|^2 + \sum_p \chi(p) \pi \log \eta + \frac{\gamma}{2\pi} m_\lambda$$

which, modulo taking the limit $\eta \to 0$, is exactly the stated result.

We may then define $f(h) = W(\nabla h, \chi) + \frac{\gamma}{2\pi} m_\lambda$, condition (ii) of Section 6.1 is then satisfied, and it is also straightforward that $f^*$ defined as in Theorem 6.1 satisfies

$$f^*(h) = \lim_{R \to \infty} \int_{K_R} f(h(\lambda + \cdot))d\lambda$$

$$= \lim_{R \to \infty} \int_{K_R} W(\nabla h, \chi * 1_{K_R}) + \frac{\gamma}{2\pi} m_\lambda = W(\nabla h) + \frac{\gamma}{2\pi} m_\lambda$$

where $W$ is now the full renormalized energy as in Definition 5.3. This allows us to use the abstract framework of Section 6.1 except there is one major assumption which is not satisfied: namely the assumption (i) that $f_\varepsilon$ must be bounded below by a constant independent of $\varepsilon$. However this assumption can easily seen not to be true! What is true is only that $f_\varepsilon$ is bounded below on average, but not pointwise. This causes one of the most serious and technical difficulties in the proof.

We point out that this problem does not occur in the “smearing out” approach to Coulomb gases, which is one of the main advantages of that approach, it does however occur in the approach of renormalizing by “cutting out holes” for the Coulomb gas. In the Ginzburg-Landau setting, we see no analogue of the smearing out method that could remedy this, in particular due to the fact that the signs of vortices can a priori be arbitrary. Instead, we rely on the ball construction lower bounds to remedy this (and this is a completely two-dimensional remedy, as is the ball construction). The lower bound stated in Theorem 8.1 do not suffice, rather we need the “improved lower bounds” that we introduced in [SS3] and which are roughly explained at the end of Section 8.1.2. Thanks to these lower bounds, which are sharp up to a constant error per vortex, we are able to show that even though the energy density associated to $f_\varepsilon$ is not bounded below pointwise, the negative part of $f_\varepsilon$ (corresponding to the subtracted vortex costs) can be “displaced” into the positive part of $f_\varepsilon$ in order to replace $f_\varepsilon$ by an equivalent energy density $g_\varepsilon$ which is pointwise bounded below, without making too much error, in the sense
that \( \|f_\varepsilon - g_\varepsilon\|_{Lip} \) is well controlled. We call this “mass displacement.” For that we have to look for energy that compensates the negative \(-\pi|\log \varepsilon'|\), and this energy is found in balls surrounded the vortices, as well as in annuli (as described at the end of Section 8.1.2) that can be up to distance \(O(1)\) away. This is done in [SS5], and at the same time it is shown that \(g_\varepsilon\) (hence \(f_\varepsilon\)) controls the number of vortices, a crucial fact whose need we mentioned above.

### 10.2.3 Upper bound

The upper bound, at least the one needed to construct a recovery sequence for minimizers, can be obtained with the same ideas as for the Coulomb gas in the proof of Proposition 6.6. There, we considered the support of the equilibrium measure and partitioned it into rectangles on which we pasted the “screened” minimizers obtained in Proposition 5.12 rescaled to have the proper density. In the Ginzburg-Landau context, this is easier since \(\mu_{\varepsilon_{x}}\) the analogue of the blown-up equilibrium measure, has a uniform density on its support. What we thus need to do is assume that this support \(\omega_{0,\varepsilon}\) is nice enough (has \(C^1\) boundary), which we can show is ensured for example by the strong requirement that \(\Omega\) be convex (we could certainly remove that condition and replace it by the assumption that \(\omega_{0,\varepsilon}\) has no cusps), and then partition it (up to a small boundary layer) into squares of size \(R \times R\). In each square we need to paste a solution of

\[
\begin{cases}
-\Delta h = 2\pi \sum_{p \in \Lambda} \delta_{p} - m_{\lambda} & \text{in } K_{R}, \\
\frac{\partial h}{\partial \nu} = 0 & \text{on } \partial K_{R},
\end{cases}
\]

with

\[
\limsup_{R \to \infty} \frac{W(\nabla h, \mathbf{1}_{K_{R}})}{|K_{R}|} \leq \min_{A_{m_{\lambda}/2\pi}} W.
\]

This is a screened minimizer of \(W\) over \(A_{m_{\lambda}/2\pi}\). The fact that such an \(h\) can be found is proven in [SS6] and follows the same outline as the proof of Proposition 5.12 in the case of \(W\), except that we do not need to first reduce to configurations with well-separated points (the screening can be accomplished for essentially generic configurations). It is however complicated by the lack of lower bound on the energy density associated to \(W(\cdot, \chi)\), which requires to go through another “mass displacement” to transform the energy density into one that is bounded below, just as we did for \(f_\varepsilon\) above. Once these screened minimizers are pasted to almost cover \(\omega_{0,\varepsilon}\), we obtain a resulting set of points (the vortices) and an associated vector field \(\nabla h\), which we extend by \(0\) outside \(\omega_{0,\varepsilon}\), and to which we add \(\nabla h_{0,\varepsilon}\). After projecting this vector field onto gradients (which can only decrease its energy, as seen in Step 5 of the proof of Proposition 6.6) and blowing down, this defines the test induced magnetic field \(h_\varepsilon\). Then there remains to build a corresponding \((u, A)\), which can be done as in the proof of the upper bound for Theorem 9.1. Because we now need the energy to be optimal at next order, we need to be more precise near the vortex cores, and plug in exactly the optimal radial profiles for vortices which give the energy \(\gamma\) per point.
10.2.4 A statement of main result

Modulo the technical difficulties mentioned above, for which we refer to [SS5, SS6], we can state the main result which holds in the setting we chose to describe. The result can be written in complete $\Gamma$-convergence form. However for simplicity, we only state it as the analogue of Theorem 6.3 together with the consequences for minimizers as in Theorem 6.4.

**Theorem 10.1** (Next order behavior of the Ginzburg-Landau functional [SS6]). Assume $\Omega$ is convex and (9.2) holds. Let $(u_\varepsilon, A_\varepsilon)$ be such that $G_\varepsilon(u_\varepsilon, A_\varepsilon) \leq G^0_\varepsilon + C h_{\text{ex}}$, and let $P_\varepsilon$ be the push-forward of the normalized Lebesgue measure on $\omega_{0,\varepsilon}$ by

$$x \mapsto \frac{1}{\sqrt{h_{\text{ex}}}} \nabla h_{\varepsilon}(x + \frac{1}{\sqrt{h_{\text{ex}}}})$$

where $h_{\varepsilon}$ is implicitly extended by 0 outside the domain $\Omega$. Then, up to extraction of a subsequence, we have $P_\varepsilon \rightarrow P$ in the weak sense of probabilities, where $P$ is some probability measure concentrated on $A_{m\lambda/2\pi}$ and

$$G_\varepsilon(u_\varepsilon, A_\varepsilon) \geq G^0_\varepsilon + h_{\text{ex}}|\omega_{0,\varepsilon}| \left( \int_{A_{m\lambda/2\pi}} W(\nabla h) dP(\nabla h) + \frac{m\lambda \gamma}{2\pi} \right) + o(h_{\text{ex}}). \quad (10.14)$$

If in addition $(u_\varepsilon, A_\varepsilon)$ minimizes $G_\varepsilon$ then $P$-a.e. $\nabla h$ minimizes $W$ over $A_{m\lambda/2\pi}$ and

$$G_\varepsilon(u_\varepsilon, A_\varepsilon) = G^0_\varepsilon + h_{\text{ex}}|\omega_{0,\varepsilon}| \left( \min_{A_{m\lambda/2\pi}} W + \frac{m\lambda \gamma}{2\pi} \right) + o(h_{\text{ex}}).$$

As announced, this provides a next order expansion of the Ginzburg-Landau energy, in a similar fashion as what we have seen for Coulomb gases with an error $o(h_{\text{ex}})$ which is also $o(1)$ per vortex. Moreover, it connects the question of minimizing Ginzburg-Landau to that of minimizing $W$. If Conjecture 1 (the conjecture that the triangular lattice minimizes $W$) was proven, then it would rigorously justify why the Abrikosov lattice appears in experiments on superconductors. The result here says that almost all the blown up configurations resemble minimizers of $W$ as $\varepsilon \rightarrow 0$. Using the method of [RNSe] (cf. end of Section 6.3) it should be possible to obtain a stronger result for all blow ups, i.e. a result of equidistribution of energy.

The assumption (9.2) has been made here for simplicity of presentation, the result in [SS6] is more general and works as long as the number of vortices diverges to infinity which happens as soon as $h_{\text{ex}} - H_{\text{c1}} \gg \log |\log \varepsilon|$, and as long as $h_{\text{ex}} \ll \frac{1}{\varepsilon^2}$. 
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