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From Vlasov–Poisson and Vlasov–Poisson–Fokker–Planck systems to incompressible Euler equations: the case with finite charge

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FROM VLASOV–POISSON AND
VLASOV–POISSON–FOKKER–PLANCK SYSTEMS TO
INCOMPRESSIBLE EULER EQUATIONS:
THE CASE WITH FINITE CHARGE

BY JULIEN BARRÉ, DAVID CHIRON, THIERRY GOUDON
& NADER MASMOUDI

ABSTRACT. — We study the asymptotic regime of strong electric fields that leads from the Vlasov–Poisson system to the Incompressible Euler equations. We also deal with the Vlasov–Poisson–Fokker–Planck system which induces dissipative effects. The originality consists in considering a situation with a finite total charge confined by a strong external field. In turn, the limiting equation is set in a bounded domain, the shape of which is determined by the external confining potential. The analysis extends to the situation where the limiting density is non-homogeneous and where the Euler equation is replaced by the Lake Equation, also called Anelastic Equation.

RÉSUMÉ (Des équations de Vlasov-Poisson et Vlasov-Poisson-Fokker-Planck aux équations d'Euler incompressibles : le cas de charge finie)

Nous étudions le régime asymptotique de forts champs électriques qui conduit du système de Vlasov-Poisson aux équations d'Euler incompressibles. Nous abordons aussi le système de Vlasov-Poisson-Fokker-Planck qui induit des effets dissipatifs additionnels. L'originalité de cette étude réside dans le fait qu'on suppose la charge totale finie et confinée par un fort champ extérieur. En conséquence, l'équation limite est posée dans un domaine borné dont la géométrie est déterminée par ce champ confinant. L'analyse s'étend au cas où la densité limite est inhomogène ; l'équation d'Euler est alors remplacée par l'équation des lacs (ou modèle anélastique).

CONTENTS

1. Introduction.....	248
2. The limit density n_e and total potential Φ_e	255
3. Asymptotic analysis.....	273
4. Asymptotic analysis of the Vlasov–Poisson–Fokker–Planck system.....	283
Appendix A. Smooth solutions of the Lake Equations.....	290
Appendix B. Construction of an extended divergence-free velocity.....	291
Appendix C. About the experimental parameters in a Magneto-Optical Trap experiment.....	292
References.....	294

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1. INTRODUCTION

1.1. THE VLASOV-POISSON EQUATION IN A CONFINING POTENTIAL. — We are interested in the behavior as ε tends to 0 of the solutions of the following Vlasov equation

$$(V) \quad \partial_t f_\varepsilon + v \cdot \nabla_x f_\varepsilon - \left(\frac{1}{\varepsilon} \nabla_x \Phi_{\text{ext}} + \nabla_x \Phi_\varepsilon \right) \cdot \nabla_v f_\varepsilon = 0,$$

where the potential Φ_ε is defined self-consistently by the Poisson equation

$$(P) \quad \Delta_x \Phi_\varepsilon = -\frac{1}{\varepsilon} \rho_\varepsilon, \quad \rho_\varepsilon(t, x) = \int f_\varepsilon(t, x, v) dv,$$

and where $\varepsilon^{-1} \Phi_{\text{ext}}$ is a strong external potential applied to the system. The problem holds in the entire space: $x \in \mathbb{R}^N$, $v \in \mathbb{R}^N$ and it is completed by an initial data with *finite* charge

$$(1) \quad f_\varepsilon|_{t=0} = f_\varepsilon^{\text{init}}, \quad \iint f_\varepsilon^{\text{init}} dv dx = \mathbf{m} \in (0, \infty).$$

Notice that Φ_ε is of size ε^{-1} and we shall consider the applied potential $\frac{1}{\varepsilon} \Phi_{\text{ext}}$ also of size ε^{-1} . The problem is motivated by the study of *non neutral plasmas* (see [14] for a review): these are collections of particles all with the same sign of charge, for instance pure electron, or pure ion plasmas. There are several methods to confine such a plasma, among which the Paul trap, which uses an oscillating electric field. The Penning trap, which uses a combination of static electric and magnetic fields, is also standard, but (V) is not directly relevant to this situation since there is no magnetic field in it. A non neutral plasma picture has also been used to describe trapped neutral atoms [37], in the regime where multiple diffusion of quasi resonant photons induces an effective interaction force between atoms which is formally similar to a Coulomb force [46]. In this case, the system is however dissipative; a standard way to take this effect into account is to add to (V) a Fokker-Planck operator acting on velocities [11]. We will also discuss this situation. In these physical examples, the small ε limit is indeed relevant in many experimental situations. Figure 1 corresponds to a numerical simulation of such an experiment. It strongly suggests the existence of a limiting fluid model where the density is nothing but the characteristic function of a ball. Our goal is to justify that, indeed, a simpler model, purely of hydrodynamic type, can be used to describe the particles in this asymptotic limit. In the appendix, we input typical orders of magnitude for a magneto-optical trap experiment and discuss the relevance of the limit in this case.

In fact, we shall see that the limiting model holds in a domain the shape of which depends on the external potential Φ_{ext} . But, to start with, we can consider a quadratic and isotropic potential, say:

$$(2) \quad \Phi_{\text{ext}}(x) = \frac{1}{2N} |x|^2,$$

where we remind the reader that N stands for the space dimension. It corresponds to the case displayed in Figure 1. The confining potential $\varepsilon^{-1} \Phi_{\text{ext}}$ tends to strongly

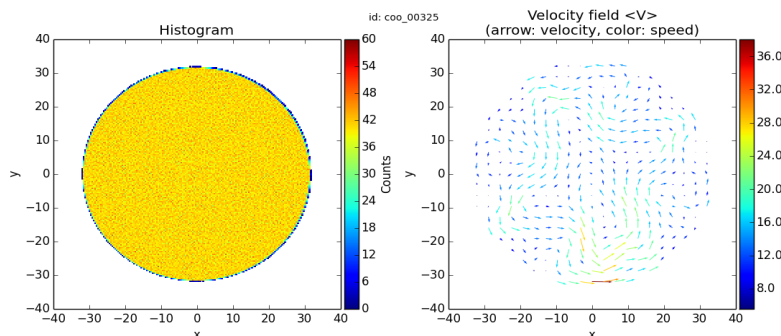


FIGURE 1. Snapshot of a 2D simulation of confined charged particles. Particles are subjected to the combination of a harmonic and isotropic external potential, a strong Coulomb repulsion, a friction and a noise. An external force has been added from the left to the right in the lower half of the cloud, in order to set the particles in motion. Left: instantaneous locally averaged density field. The density is almost uniform inside a ball, and almost zero outside. Right: instantaneous locally averaged velocity field. (By courtesy of A. Olivetti [40].)

localize in space the particle density. On the support of the limiting density ρ , the electric force $\varepsilon^{-1}\nabla_x\Phi_{\text{ext}} + \nabla_x\Phi_\varepsilon$ should be of order one. By (P), this imposes that

$$\Delta\Phi_{\text{ext}} + \varepsilon\Delta\Phi_\varepsilon = \varepsilon\nabla_x \cdot (\varepsilon^{-1}\nabla_x\Phi_{\text{ext}} + \nabla_x\Phi_\varepsilon) = \Delta\Phi_{\text{ext}} - \rho_\varepsilon = 1 - \rho_\varepsilon$$

is of order $\mathcal{O}(\varepsilon)$ on the support of ρ for the potential (2). Clearly, due to the condition of finite charge (1), the limiting density cannot be constant uniformly on the whole space. The intuition is that the limiting density has the same radial symmetries as both the external potential (2) and the Poisson kernel, see (14) below. Actually, we shall prove some convergence of ρ_ε to

$$(3) \quad n_e(x) = \mathbf{1}_{B(0,R)}(x),$$

where $\mathbf{1}_U$ denotes the characteristic function of the set U . The radius R depends on the total mass \mathbf{m} so that the charge constraint (1) is fulfilled. In order to find a hydrodynamic description of the particles, it is convenient to associate to the particle distribution function f_ε the following macroscopic quantities

$$\begin{aligned} \text{Current:} \quad J_\varepsilon(t, x) &:= \int v f_\varepsilon(t, x, v) dv, \\ \text{Momentum flux:} \quad \mathbb{P}_\varepsilon(t, x) &:= \int v \otimes v f_\varepsilon(t, x, v) dv. \end{aligned}$$

It turns out that the current looks like

$$(4) \quad J_\varepsilon(t, x) = \rho_\varepsilon(t, x)V_\varepsilon(t, x) \xrightarrow{\varepsilon \rightarrow 0} n_e(x)V(t, x) = \mathbf{1}_{B(0,R)}(x)V(t, x),$$

where V solves the Incompressible Euler system in $B(0, R)$:

$$(IE) \quad \begin{cases} \partial_t V + \nabla_x \cdot (V \otimes V) + \nabla_x p = 0, \\ \nabla_x \cdot V = 0, \end{cases}$$

with an appropriate initial condition, and no-flux boundary condition on $\partial B(0, R)$. In (IE), the pressure p appears as the Lagrange multiplier associated with the constraint that V is divergence free. This incompressibility condition comes from charge conservation: integrating (V) with respect to the velocity variable v , we get

$$(5) \quad \partial_t \rho_\varepsilon + \nabla_x \cdot J_\varepsilon = 0.$$

Letting ε go to 0, with (3) and (4), we deduce that V is solenoidal. Obtaining the evolution equation for V is more intricate.

The analysis of such asymptotic problems goes back to [5], where a specific modulated energy method was introduced. It has been revisited in [35], still by using a modulated energy method, but which is able to account for oscillations present within the system. Accordingly, more general initial data can be dealt with in [35]. However, these results hold either on the torus \mathbb{T}^N , or in the whole space with data having infinite charge, that is $\iint f(x, v) dv dx = \infty$. A case with finite charge, but a different Poisson equation which leads to a compressible hydrodynamic limit, has been considered in [23], again with a modulated energy. Further results on the quasi-neutral regime can be found in [21, 22, 23, 24, 25]. Our goal in this article is twofold:

- To prove the convergence to (IE) in the case of a trapped system, with finite charge. Even though our proof also relies on a modulated energy functional, there are new difficulties: the shape of the domain on which the limiting equation (IE) holds is determined by the external potential Φ_{ext} , and a careful treatment of the boundary is needed.

- To prove the convergence to the analog of (IE) in the case of a trapped dissipative system.

Both improvements are relevant for experiments on non neutral plasmas or large magneto-optical traps.

1.2. STATEMENT OF THE RESULTS. — In what follows we shall deal with a smooth solution $(t, x) \mapsto V(t, x) \in \mathbb{R}^N$ (possibly defined on a small enough time interval $[0, T]$) of the incompressible Euler equation (IE) set on the ball $B(0, R)$, completed with no-flux boundary condition

$$(6) \quad V(t, x) \cdot \nu(x)|_{|x|=R} = 0,$$

where $\nu(x)$ denotes the outward unit vector at $x \in \partial B(0, R)$ (namely $\nu(x) = x/|x|$). We work with solutions V that belongs to $L^\infty(0, T; H^s(B(0, R)))$, for a certain $s > 0$ large enough.

THEOREM 1.1 ([44, 45]). — *Let $V^{\text{init}} : B(0, R) \rightarrow \mathbb{R}^N$ be a divergence free vector field in H^s , with $s > 1 + N/2$, satisfying the no-flux condition $V^{\text{init}} \cdot \nu = 0$ on $\partial B(0, R)$.*

There exists $T > 0$ and a unique solution $V \in L^\infty(0, T; H^s(B(0, R)))$ of (IE) with the no-flux condition (6). Moreover, we have

$$\sup_{0 \leq t \leq T} \left(\|V(t)\|_{H^s} + \|\partial_t V(t)\|_{H^{s-1}} + \|\nabla_x p(t)\|_{H^s} + \|\partial_t \nabla_x p(t)\|_{H^{s-1}} \right) \leq C(T)$$

for some positive constant $C(T)$ depending on T and the initial datum.

If $N = 1$, the only divergence free vector field V^{init} satisfying (6) is $V^{\text{init}} \equiv 0$ and then the solution given in Theorem 1.1 is $V \equiv 0$.

For further purposes, we need to consider an extension \mathcal{V} of the solution V to (IE) with (6), defined on the whole space and compactly supported. Namely we require $\mathcal{V} \in L^\infty(0, T; H^s(\mathbb{R}^N))$ to satisfy

$$(7) \quad \mathcal{V}|_{B(0,R)} = V, \quad \mathcal{V}|_{\mathbb{R}^N \setminus B(0,2R)} = 0, \quad \mathcal{V}(t, x) \cdot \nu(x)|_{|x|=R} = 0.$$

For the construction of such an extension, we refer to [33, Chap. I: Th. 2.1 p. 17 & Th. 8.1 p. 42]. For an extension which is in addition divergence-free, see Lemma B.1 in the appendix.

In order to state our first result, we need to introduce an auxiliary potential function Φ_e . Suppose that (3) indeed holds true. Then, by using (P) and $\Delta \Phi_{\text{ext}} = 1$ for the potential (2), we infer, for $\varepsilon \rightarrow 0$,

$$\Delta(\Phi_{\text{ext}} + \varepsilon \Phi_\varepsilon) = \Delta \Phi_{\text{ext}} - \rho_\varepsilon \longrightarrow \Delta \Phi_{\text{ext}} - \mathbf{1}_{B(0,R)} = \mathbf{1}_{\mathbb{R}^N \setminus B(0,R)}.$$

Moreover, since we want the electric force $\varepsilon^{-1} \nabla_x \Phi_{\text{ext}} + \nabla_x \Phi_\varepsilon = \varepsilon^{-1} (\nabla_x \Phi_{\text{ext}} + \varepsilon \nabla_x \Phi_\varepsilon)$ to be of order one on the ball $B(0, R)$, this imposes $\Phi_{\text{ext}} + \varepsilon \Phi_\varepsilon$ to be close to a constant, say zero, on the ball $B(0, R)$. It is therefore natural to look for a solution Φ_e to the Poisson problem

$$(8) \quad \Delta \Phi_e(x) = 1 - n_e(x) = \mathbf{1}_{\mathbb{R}^N \setminus B(0,R)}, \quad \Phi_e = 0 \text{ in } B(0, R).$$

In this specific case, we can find an explicit radially symmetric solution:

$$(9) \quad \Phi_e(x) = \mathbf{1}_{\mathbb{R}^N \setminus B(0,R)} \times \begin{cases} \frac{|x|^2}{2N} + \frac{R^N}{N(N-2)|x|^{N-2}} - \frac{R^2}{2(N-2)} & \text{if } N > 2, \\ \frac{|x|^2 - R^2}{4} - \frac{R^2}{2} \ln(|x|/R) & \text{if } N = 2, \\ \frac{1}{2}(|x| - R)^2 & \text{if } N = 1. \end{cases}$$

With Φ_e and n_e in hand, we split the Poisson equation (P) as follows, where n_e is defined in (3),

$$\begin{aligned} \Delta_x \Phi_\varepsilon(t, x) &= \frac{1 - n_e(x)}{\varepsilon} + \frac{n_e(x) - \rho_\varepsilon(t, x)}{\varepsilon} - \frac{1}{\varepsilon} \Delta \Phi_{\text{ext}} \\ &= \frac{1}{\varepsilon} \Delta_x \Phi_e(x) + \frac{1}{\sqrt{\varepsilon}} \Delta_x \Psi_\varepsilon(t, x) - \frac{1}{\varepsilon} \Delta \Phi_{\text{ext}}, \end{aligned}$$

namely, we have

$$(10) \quad \Phi_\varepsilon(x) + \frac{1}{\varepsilon} \Phi_{\text{ext}} = \frac{1}{\varepsilon} \Phi_e(x) + \frac{1}{\sqrt{\varepsilon}} \Psi_\varepsilon(t, x), \quad \Delta_x \Psi_\varepsilon(t, x) = \frac{1}{\sqrt{\varepsilon}} (n_e(x) - \rho_\varepsilon(t, x)),$$

where Ψ_ε represents the fluctuations of the potential. According to [5], we introduce a modulated energy:

$$\mathcal{H}_{\mathcal{V},\varepsilon} := \frac{1}{2} \iint |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon \, dv \, dx.$$

When the external potential is given by (2), we shall establish the following statement.⁽¹⁾

THEOREM 1.2. — *Set*

$$R = \left(\frac{\mathbf{m}}{|B(0,1)|} \right)^{1/N}.$$

Let $V^{\text{init}} \in H^s(B(0,R))$ satisfy $\nabla_x \cdot V^{\text{init}} = 0$ and the no-flux condition (6). Denote by V the solution, on $[0,T]$, to (IE) with the no-flux condition (6) given in Theorem 1.1. Consider \mathcal{V} a smooth extension of V satisfying the conditions (7). Let $f_\varepsilon^{\text{init}} : \mathbb{R}^N \times \mathbb{R}^N \rightarrow [0, \infty)$ be a sequence of integrable functions that satisfy the following requirements

$$(11) \quad \left\{ \begin{array}{l} \iint f_\varepsilon^{\text{init}} \, dv \, dx = \mathbf{m}, \\ \lim_{\varepsilon \rightarrow 0} \left\{ \frac{1}{2} \iint |v - \mathcal{V}^{\text{init}}|^2 f_\varepsilon^{\text{init}} \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon^{\text{init}}|^2 \, dx \right. \\ \left. + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon^{\text{init}} \, dv \, dx \right\} = 0. \end{array} \right.$$

Then, the associated solution f_ε of the Vlasov–Poisson equation (V)–(P) satisfies, as $\varepsilon \rightarrow 0$,

- (i) ρ_ε converges to n_e in $C^0([0,T]; \mathcal{M}^1(\mathbb{R}^N))$ -weak- \star ;
- (ii) $\mathcal{H}_{\mathcal{V},\varepsilon}$ converges to 0 uniformly on $[0,T]$;
- (iii) J_ε converges to J in $\mathcal{M}^1([0,T] \times \mathbb{R}^N)$ weakly- \star , the limit J lies in $L^\infty(0,T; L^2(\mathbb{R}^N))$ and satisfies $J|_{[0,T] \times B(0,R)} = V$, $\nabla_x \cdot J = 0$ and $J \cdot \nu(x)|_{\partial B(0,R)} = 0$.

REMARK 1.3

(i) Here, we were not very precise about the type of solutions to the Vlasov–Poisson system (V)–(P) we are considering. We refer to [41, 34] for the construction of global regular solutions to the system and some extra conditions to ensure the propagation of regularity. There are also weaker notions of solutions (weak solutions or renormalized solutions) to which our theorem can apply. We refer the reader to the introduction of [34] for a discussion about these solutions.

(ii) The second part of the hypothesis (11) imposes that the initial modulated energy is small; this is a strong hypothesis on the initial data. As pointed out in [5], it corresponds to assume that the particles’ temperature vanishes: it reduces the spreading of the particles distribution in velocity. When the problem is set on the torus, or on the whole space with infinite charge, it can be relaxed, see [35], but still for

⁽¹⁾Throughout the paper, we denote by $\mathcal{M}^1(X)$ the space of bounded measures on $X \subset \mathbb{R}^D$. It identifies with the dual space of the separable space $C_0^0(X)$ of the continuous functions that vanish at infinity.

a “cold particles” regime, though. In the present framework, going beyond (11) would certainly require a fine description of boundary layers on $\{|x| = R\}$. Assuming (11), Point (ii) of the theorem then ensures that the modulated energy remains small at later times. As typical initial data satisfying (11), we can take

$$f_\varepsilon^{\text{init}}(x, v) = \frac{n_e(x) - \delta_\varepsilon \Delta \chi(x)}{\sigma_\varepsilon^N} G\left(\frac{v - \mathcal{V}^{\text{init}}(x)}{\sigma_\varepsilon}\right),$$

where $\chi \in C_c^\infty(B(0, R))$ and where G is a nonnegative function that belongs to the Schwartz space and satisfies $\int G \, dv = 1$ (for instance, G is a normalized Gaussian $G(v) = (2\pi)^{-N/2} \exp(-|v|^2/2)$). Then, we choose $\sigma_\varepsilon \rightarrow 0$ as $\varepsilon \rightarrow 0$ and $\delta_\varepsilon = o(\sqrt{\varepsilon})$ (so that $n_e - \delta_\varepsilon \Delta \chi \geq 0$ for ε small enough). Indeed, we easily obtain

$$\begin{aligned} \iint \Phi_e f_\varepsilon^{\text{init}} \, dv \, dx &= 0, \\ \iint |v - \mathcal{V}^{\text{init}}|^2 f_\varepsilon^{\text{init}} \, dv \, dx &= \sigma_\varepsilon^2 \left(\int |v|^2 G(v) \, dv \right) \left(\int n_e \, dx \right) \rightarrow 0 \end{aligned}$$

and
$$\Psi_\varepsilon = \frac{\delta_\varepsilon}{\sqrt{\varepsilon}} \chi,$$

hence
$$\int |\nabla_x \Psi_\varepsilon^{\text{init}}|^2 \, dx = \frac{\delta_\varepsilon^2}{\varepsilon} \int |\nabla_x \chi|^2 \, dx \rightarrow 0.$$

We wish to extend this analysis by dealing with more general external potentials. We distinguish two situations depending on the expression of the external potential:

- The quadratic potential

$$(12) \quad \Phi_{\text{ext}}(x) = \frac{1}{2} \sum_{j=1}^N \frac{x_j^2}{\lambda_j^2},$$

with $\lambda_j > 0$, $1 \leq j \leq N$, in dimension $N \geq 2$ is typical to model non neutral plasmas [14] or magneto-optical traps experiments. In this case $\Delta \Phi_{\text{ext}}$ is still a constant, that therefore determines the value of the (uniform) particle density n_e on its support. But the problem has lost its symmetries and the shape of the support becomes non trivial. We shall see that ρ_ε tends to a uniform distribution n_e , supported in an ellipsoid. However, we point out that the support of n_e does not coincide with a level set of Φ_{ext} . An example with $N = 2$ is given in Figure 2. The potential Φ_e can be computed rather explicitly, and Theorem 1.2 generalizes directly. See Section 2.1 for a precise statement.

– In the case of a non quadratic potential, under suitable hypotheses on Φ_{ext} , the limiting density n_e still has a compact support \mathcal{K} and is still given on \mathcal{K} by $n_e = \Delta \Phi_{\text{ext}}$. However, n_e is clearly no longer constant on \mathcal{K} . The identification of \mathcal{K} and n_e relies on variational techniques, with connection to the obstacle problem. It is still possible to prove the analog of Theorem 1.2, but, since n_e becomes non homogeneous, instead of (IE) the limiting equations are now the so-called Lake

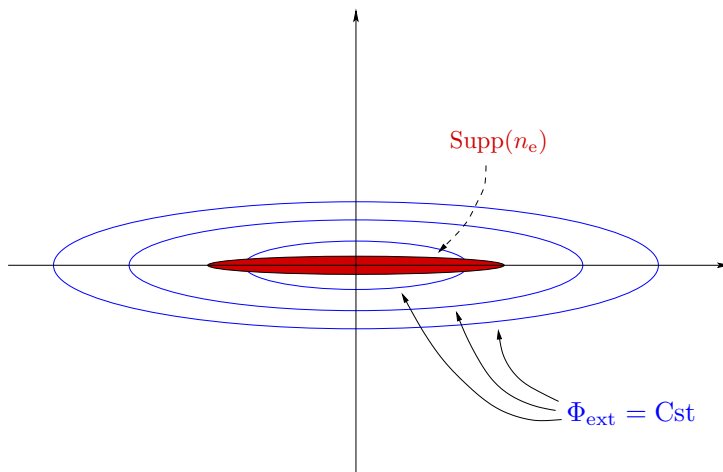


FIGURE 2. Some level sets of Φ_{ext} and the support of n_e .

Equations, see e.g. [32]:

$$(LE) \quad \begin{cases} \partial_t V + V \cdot \nabla_x V + \nabla_x p = 0, \\ \nabla_x \cdot (n_e V) = 0. \end{cases}$$

Such model — also referred to as the Anelastic Equations — arise in the modelling of atmospheric flows [38]; we refer the reader to [36] for the justification of a derivation from the compressible Navier-Stokes system. As a matter of fact, we can observe that the first equation in (LE) may be written in the following conservative form $\partial_t(n_e V) + \nabla_x \cdot (n_e V \otimes V) + n_e \nabla_x p = 0$. The construction of Φ_e and \mathcal{K} , and a precise statement of the corresponding convergence theorem can be found in Section 2.2.

Motivated by actual experiments, we will also generalize the results to the case where a Fokker-Planck operator acting on velocities is added to Equation (V). Our starting point then becomes:

$$(VFP) \quad \partial_t f_\varepsilon + v \cdot \nabla_x f_\varepsilon - \nabla_x \Phi_\varepsilon \cdot \nabla_v f_\varepsilon = L f_\varepsilon,$$

with

$$L f = \nabla_v \cdot (v f + \theta \nabla_v f) = \theta \nabla_v \cdot (M_{0,\theta} \nabla_v (f/M_{0,\theta})), \quad M_{0,\theta}(v) = \frac{1}{(2\pi\theta)^{N/2}} e^{-|v|^2/(2\theta)},$$

for some $\theta > 0$. Equation (VFP) is still coupled to the Poisson equation (P). Using a modified modulated energy, we are able to show in this case that solutions f_ε of (VFP) and (P) with well-prepared data converge when ε and θ tends to 0 in the sense of Theorem 1.2 to $n_e V$, where V is now the solution of the Lake Equation with friction

$$(13) \quad \begin{cases} \partial_t V + V \cdot \nabla_x V + \nabla_x p + V = 0, \\ \nabla_x \cdot (n_e V) = 0. \end{cases}$$

On the boundary, we still have the no-flux condition (6). For the sake of completeness, the necessary analog of Theorem 1.1 for the systems (LE) and (13) is sketched in appendix A. See Section 4 for a precise statement on the asymptotic behavior of (VFP) and its proof.

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2. THE LIMIT DENSITY n_e AND TOTAL POTENTIAL Φ_e

As said above, we have a clear intuition and explicit formulas for the equilibrium distribution n_e and the potential Φ_e in the specific case of the isotropic external potential (2). Let us discuss in further details how Φ_{ext} determines n_e and its support, and how the auxiliary potential Φ_e , which plays a crucial role in the analysis through the decomposition (10), can be defined.

We remind the reader the definition of the fundamental solution, hereafter denoted Γ , of $(-\Delta)$ (mind the sign) in the whole space \mathbb{R}^N :

$$(14) \quad \Gamma(x) := \begin{cases} \frac{1}{N(N-2)|B(0,1)| \cdot |x|^{N-2}} & \text{if } N > 2, \\ -\ln|x|/2\pi & \text{if } N = 2, \\ -|x|/2 & \text{if } N = 1. \end{cases}$$

2.1. THE CASE OF A GENERAL QUADRATIC POTENTIAL. — Let us consider in this section the case of a quadratic potential (12). We have $\Delta\Phi_{\text{ext}} = \sum_{j=1}^N 1/\lambda_j^2 > 0$ which is constant in space. It gives the value of the equilibrium density on its support since we still expect $\rho_\varepsilon \rightarrow \mathbf{1}_{\mathcal{H}} \Delta\Phi_{\text{ext}}$. But it remains to determine this support $\text{Supp}(n_e) = \mathcal{H} \subset \mathbb{R}^N$ on which we have the volume constraint

$$\mathfrak{m} = \int n_e \, dx = |\mathcal{H}| \sum_{j=1}^N \frac{1}{\lambda_j^2}$$

coming from (1). Note that a quick computation reveals that \mathcal{H} can be neither radially symmetric, nor a level set of Φ_{ext} unless $\lambda_1 = \dots = \lambda_N$.

In order to extend Theorem 1.2 for a potential as in (12), we need to construct a domain $\mathcal{H} \subset \mathbb{R}^N$ and a function $\Phi_e : \mathbb{R}^N \rightarrow \mathbb{R}$ such that

$$(15) \quad \Delta\Phi_e(x) = \left(\sum_{j=1}^N \frac{1}{\lambda_j^2} \right) \mathbf{1}_{\mathbb{R}^N \setminus \mathcal{H}}, \quad \Phi_e = 0 \quad \text{in } \mathcal{H}.$$

The starting point is the observation that given $a = (a_1, \dots, a_N) \in (\mathbb{R}_+^*)^N$, then the characteristic function of the ellipsoid

$$\mathcal{K}_a = \left\{ x \in \mathbb{R}^N; \sum_{j=1}^N x_j^2/a_j^2 \leq 1 \right\}$$

generates an electric potential which is *quadratic* inside the ellipsoid. This can be found for instance in [28, Chap. VII, §6]; the computation there is for $N = 3$, but the extension to the case $N \geq 3$ is straightforward, and the two-dimensional case is treated by using arguments from complex analysis in [19].

For $x \in \mathbb{R}^N$, we denote by $\sigma_a(x)$ the largest solution of the equation

$$\sum_{j=1}^N \frac{x_j^2}{a_j^2 + \varsigma} = 1$$

(with $\varsigma \in \mathbb{R}$ as unknown). Consequently, $x \in \mathcal{K}_a$ holds if and only if $\sigma_a(x) \leq 0$. By convention, $\sigma_a(0) = -\infty$. This quantity can be seen as an equivalent of the radial coordinate in the ellipsoidal coordinate system. It allows us to construct a solution to (15) where \mathcal{K} is an ellipsoid, the coefficients of which depend on the mass \mathbf{m} and the λ_j 's.

PROPOSITION 2.1. — *Let $a = (a_1, \dots, a_N) \in (\mathbb{R}_+^*)^N$.*

(i) [28] *If $N \geq 3$, then*

$$\Gamma \star \mathbf{1}_{\mathcal{K}_a}(x) = \frac{1}{4} \left(\prod_{j=1}^N a_j \right) \times \begin{cases} \int_{\sigma_a(x)}^{+\infty} \left(1 - \sum_{j=1}^N \frac{x_j^2}{a_j^2 + s} \right) \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds & \text{if } \sigma_a(x) \geq 0, \\ \int_0^{+\infty} \left(1 - \sum_{j=1}^N \frac{x_j^2}{a_j^2 + s} \right) \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds & \text{if } \sigma_a(x) \leq 0. \end{cases}$$

(ii) *If $N = 2$, then*

$$\Gamma \star \mathbf{1}_{\mathcal{K}_a}(x) = \frac{1}{4} (a_1 a_2) \times \begin{cases} -\ln \left(\sigma_a(x) + \frac{a_1^2 + a_2^2}{2} + \sqrt{(a_1^2 + \sigma_a(x))(a_2^2 + \sigma_a(x))} \right) \\ \quad - \int_{\sigma_a(x)}^{+\infty} \sum_{j=1}^2 \frac{x_j^2}{a_j^2 + s} \frac{ds}{\sqrt{(a_1^2 + s)(a_2^2 + s)}} & \text{if } \sigma_a(x) \geq 0, \\ -\ln \left(\frac{1}{2} (a_1 + a_2)^2 \right) - \int_0^{+\infty} \sum_{j=1}^2 \frac{x_j^2}{a_j^2 + s} \frac{ds}{\sqrt{(a_1^2 + s)(a_2^2 + s)}} & \text{if } \sigma_a(x) \leq 0. \end{cases}$$

REMARK 2.2. — An alternative point of view for the two dimensional case is to work with the electric field instead of the potential. We refer to [19] for expressions of the electric field generated by ellipses in $N = 2$. In the case of a uniform charge distribution, the electric field is linear inside the ellipse, with the same coefficients for the quadratic terms as those coming from the expression in (ii).

We define the mapping $\mathcal{Z} : (\mathbb{R}_+^*)^N \rightarrow (\mathbb{R}_+^*)^N$ by

$$(16) \quad \forall k \in \{1, \dots, N\}, \quad \mathcal{Z}_k(\alpha) = \int_0^{+\infty} \frac{1}{\alpha_k + s} \left(\prod_{j=1}^N (\alpha_j + s) \right)^{-1/2} ds > 0.$$

From Proposition 2.1, we know that the potential generated by $\mathbf{1}_{\mathcal{K}_a}$ is quadratic inside \mathcal{K}_a , up to an additive constant. The coefficients of the quadratic terms are the $-(\prod_{j=1}^N a_j)\mathcal{Z}_k(\alpha)/4$, $1 \leq k \leq N$. The idea to make the connexion with the external potential Φ_{ext} is now to adapt the a_j 's so that the quadratic terms in $\Gamma \star \mathbf{1}_{\mathcal{K}_a}$ (inside \mathcal{K}_a) cancel out the quadratic terms of Φ_{ext} , so that $(\Delta\Phi_{\text{ext}})\Gamma \star \mathbf{1}_{\mathcal{K}_a} + \Phi_{\text{ext}}$ is constant in \mathcal{K}_a . We observe that $\prod_{j=1}^N a_j$ is related to the total charge of the ellipsoid \mathcal{K}_a since

$$\mathbf{m} = \int n_e = |\mathcal{K}_a| \sum_{j=1}^N \lambda_j^{-2} = |B(0, 1)| \left(\prod_{j=1}^N a_j \right) \sum_{j=1}^N \lambda_j^{-2}.$$

We shall thus need to solve equations in a of the form $\mathcal{Z}(a_1^2, \dots, a_N^2) = z$, where $z \in (\mathbb{R}_+^*)^N$ is given. Therefore, we are interested in showing that $\mathcal{Z} : (\mathbb{R}_+^*)^N \rightarrow (\mathbb{R}_+^*)^N$ is a smooth diffeomorphism. When $N = 2$, explicit computations may be carried out.

PROPOSITION 2.3. — *Assume $N = 2$. Then, for any $\alpha \in (\mathbb{R}_+^*)^2$*

$$\mathcal{Z}(\alpha) = (\mathcal{Z}_1(\alpha), \mathcal{Z}_2(\alpha)) = \left(\frac{2}{\alpha_1 + \sqrt{\alpha_1\alpha_2}}, \frac{2}{\alpha_2 + \sqrt{\alpha_1\alpha_2}} \right).$$

Moreover, $\mathcal{Z} : (\mathbb{R}_+^)^2 \rightarrow (\mathbb{R}_+^*)^2$ is a smooth diffeomorphism and its inverse is given by*

$$\mathcal{Z}^{-1}(z) = ((\mathcal{Z}^{-1})_1(z), (\mathcal{Z}^{-1})_2(z)) = \left(\frac{2z_2}{z_1(z_1 + z_2)}, \frac{2z_1}{z_2(z_1 + z_2)} \right).$$

Proof. — The explicit formula for $\mathcal{Z}(\alpha)$ comes by computing the Abelian integral

$$\int_0^{+\infty} \frac{ds}{(\alpha_1 + s)^{3/2}(\alpha_2 + s)^{1/2}} = \int_0^{+\infty} \frac{d}{ds} \left(\frac{2}{\alpha_2 - \alpha_1} \sqrt{\frac{\alpha_1 + s}{\alpha_2 + s}} \right) ds = \frac{2}{\alpha_1 + \sqrt{\alpha_1\alpha_2}},$$

for $\alpha_1 \neq \alpha_2$, and the formula holds true when $\alpha_1 = \alpha_2$ as well. The formula for the inverse then follows by direct substitution. □

For $N \geq 3$, we no longer have simple expressions for \mathcal{Z} . However, we shall prove that $\mathcal{Z} : (\mathbb{R}_+^*)^N \rightarrow (\mathbb{R}_+^*)^N$ is a smooth diffeomorphism by using the fact that $\mathcal{Z} : (\mathbb{R}_+^*)^N \rightarrow (\mathbb{R}_+^*)^N$ is a gradient vector field associated with a strictly concave function.

PROPOSITION 2.4. — *Assume $N \geq 2$ and let us define the function $\zeta : (\mathbb{R}_+^*)^N \rightarrow \mathbb{R}$ by:*

$$\zeta(\alpha) = \begin{cases} - \int_0^{+\infty} \left(\prod_{k=1}^N (\alpha_k + s) \right)^{-1/2} ds & \text{if } N \geq 3, \\ 4 \ln(\sqrt{\alpha_1} + \sqrt{\alpha_2}) & \text{if } N = 2. \end{cases}$$

Then, $\zeta : (\mathbb{R}_+^*)^N \rightarrow \mathbb{R}$ is smooth, strictly concave and it satisfies $\nabla\zeta = \mathcal{Z}$. Furthermore, $\nabla\zeta = \mathcal{Z} : (\mathbb{R}_+^*)^N \rightarrow (\mathbb{R}_+^*)^N$ is a smooth diffeomorphism and for any $z \in (\mathbb{R}_+^*)^N$, $\mathcal{Z}^{-1}(z)$ is the unique minimizer for

$$(17) \quad \inf_{\alpha \in (\mathbb{R}_+^*)^N} (z \cdot \alpha - \zeta(\alpha)).$$

In (17) we recognize the minimization problem that defines the Legendre transform of ζ . This gives a way to compute numerically $\mathcal{Z}^{-1}(z)$ through the minimization of a convex function.

Proof. — The smoothness of ζ is clear and $\nabla\zeta = \mathcal{Z}$ follows from direct computations. If $N = 2$, the strict concavity of ζ is straightforward and the fact that $\nabla\zeta = \mathcal{Z} : (\mathbb{R}_+^*)^2 \rightarrow (\mathbb{R}_+^*)^2$ is a smooth diffeomorphism comes from Proposition 2.3: for any $z \in (\mathbb{R}_+^*)^2$, $\mathcal{Z}^{-1}(z)$ is a critical point of the strictly convex (since ζ is strictly concave) function $\alpha \mapsto z \cdot \alpha - \zeta(\alpha)$, hence is the unique minimizer of that function. We assume now $N \geq 3$. Then, for each $s \in \mathbb{R}_+$, the function

$$\varpi_s : \alpha \in (\mathbb{R}_+^*)^N \mapsto \left(\prod_{k=1}^N (\alpha_k + s) \right)^{-1/2}$$

is logarithmically strictly convex since $\ln \circ \varpi_s(\alpha) = (-1/2) \sum_{k=1}^N \ln(\alpha_k + s)$ and

$$\text{Hess}(\ln \circ \varpi_s, \alpha) = (1/2) \text{Diag}((\alpha_1 + s)^{-2}, \dots, (\alpha_N + s)^{-2}).$$

Consequently, $-\zeta(\alpha) = \int_0^{+\infty} \varpi_s(\alpha) ds$ is a strictly convex function of α . Let us show that the Jacobian determinant of \mathcal{Z} never vanishes, that is $\text{Hess}(\zeta, \alpha)$ is everywhere negative definite. For that purpose, for $v \in \mathbb{R}^N$, we write

$$-v^T \text{Hess}(\zeta, \alpha)v = \int_0^{+\infty} v^T \text{Hess}(\varpi_s, \alpha)v ds,$$

and thus it suffices to show that $\text{Hess}(\varpi_s, \alpha)$ is positive definite for any $s \geq 0$. Now, we write $\varpi_s(\alpha) = \exp(\ln \circ \varpi_s(\alpha))$, thus

$$\partial_{j,k}^2 \varpi_s(\alpha) = \exp(\ln \circ \varpi_s(\alpha)) [\partial_{j,k}^2 (\ln \circ \varpi_s)(\alpha) + \partial_j (\ln \circ \varpi_s)(\alpha) \partial_k (\ln \circ \varpi_s)(\alpha)].$$

Therefore, if $v \neq 0$, we obtain

$$\begin{aligned} v^T \text{Hess}(\varpi_s, \alpha)v &= \varpi_s(\alpha) \left[v^T \text{Hess}(\ln \circ \varpi_s, \alpha)v + \left(\sum_{j=1}^N v_j \partial_j (\ln \circ \varpi_s)(\alpha) \right)^2 \right] \\ &\geq \varpi_s(\alpha) v^T \text{Hess}(\ln \circ \varpi_s, \alpha)v = \varpi_s(\alpha) \sum_{j=1}^N \frac{v_j^2}{2(\alpha_j + s)^2} > 0, \end{aligned}$$

as wished.

Let us now fix $z \in (\mathbb{R}_+^*)^N$ and consider the minimization problem (17). In view of the negativity of ζ , this infimum μ belongs to $[0, +\infty)$. Since ζ is strictly concave, this problem has at most one minimizer. Let us show that it has at least one by considering a minimizing sequence $(\alpha^n)_{n \geq 0} \in (\mathbb{R}_+^*)^N$. We claim that the sequence $(\alpha^n)_{n \geq 0}$ is bounded. Indeed, we have $z \cdot \alpha^n - \zeta(\alpha^n) \rightarrow \mu \in \mathbb{R}_+$, and since $\zeta \leq 0$,

this implies $z \cdot \alpha^n = \zeta(\alpha^n) + \mu + o(1) \leq \mu + o(1)$. Using that all components of z are positive, the claim follows. As a consequence, we may assume, up to a subsequence, that there exists $\alpha = (\alpha_1, \dots, \alpha_N) \in \mathbb{R}_+^N$ such that $\alpha^n \rightarrow \alpha$ as $n \rightarrow +\infty$. In particular, $\zeta(\alpha^n) = z \cdot \alpha^n - \mu + o(1)$ converges. We now prove that at most two components of α vanish. For otherwise, Fatou's lemma would yield

$$\begin{aligned} +\infty &= \int_0^{+\infty} \left(\prod_{k=1}^N (\alpha_k + s) \right)^{-1/2} ds = \int_0^{+\infty} \liminf_{n \rightarrow +\infty} \left(\prod_{k=1}^N (\alpha_k^n + s) \right)^{-1/2} ds \\ &\leq \liminf_{n \rightarrow +\infty} \int_0^{+\infty} \left(\prod_{k=1}^N (\alpha_k^n + s) \right)^{-1/2} ds = \liminf_{n \rightarrow +\infty} (-\zeta(\alpha^n)), \end{aligned}$$

contradicting the convergence of $(\zeta(\alpha^n))_{n \in \mathbb{N}}$. It remains to show that α has no zero component to ensure that $\mu + o(1) = z \cdot \alpha^n - \zeta(\alpha^n) \rightarrow z \cdot \alpha - \zeta(\alpha)$ so that $\alpha \in (\mathbb{R}_+^*)^N$ is actually a minimizer for (17). We then assume that $\alpha_1 = 0$, for instance, and show that for sufficiently small $\delta > 0$,

$$z \cdot (\delta, \alpha_2, \dots, \alpha_N) - \zeta(\delta, \alpha_2, \dots, \alpha_N) < z \cdot (0, \alpha_2, \dots, \alpha_N) - \zeta(0, \alpha_2, \dots, \alpha_N).$$

This reaches a contradiction for n large enough. We thus compute

$$\begin{aligned} D(\delta) &= (z \cdot (\delta, \alpha_2, \dots, \alpha_N) - \zeta(\delta, \alpha_2, \dots, \alpha_N)) \\ &\quad - (z \cdot (0, \alpha_2, \dots, \alpha_N) - \zeta(0, \alpha_2, \dots, \alpha_N)) \\ &= z_1 \delta + \int_0^{+\infty} (\delta + s)^{-1/2} \left(\prod_{k=2}^N (\alpha_k + s) \right)^{-1/2} ds - \int_0^{+\infty} s^{-1/2} \left(\prod_{k=2}^N (\alpha_k + s) \right)^{-1/2} ds \\ &= \delta \left(z_1 - \int_0^{+\infty} \frac{1}{s^{1/2}(\delta + s)^{1/2}[s^{1/2} + (\delta + s)^{1/2}]} \left(\prod_{k=2}^N (\alpha_k + s) \right)^{-1/2} ds \right). \end{aligned}$$

As $\delta \rightarrow 0$, we have, by monotone convergence,

$$\begin{aligned} \int_0^{+\infty} \frac{\left(\prod_{k=2}^N (\alpha_k + s) \right)^{-1/2}}{s^{1/2}(\delta + s)^{1/2}[s^{1/2} + (\delta + s)^{1/2}]} ds \\ \longrightarrow \int_0^{+\infty} \frac{1}{2s^{3/2}} \left(\prod_{k=2}^N (\alpha_k + s) \right)^{-1/2} ds = +\infty, \end{aligned}$$

hence for δ sufficiently small, $D(\delta) < 0$, as claimed. Therefore, $\alpha \in (\mathbb{R}_+^*)^N$ and α is a minimizer for (17). It then follows that $\nabla \zeta(\alpha) = z$ as wished. \square

We may now construct a solution to (15).

COROLLARY 2.5 (Construction of the function Φ_e for quadratic potentials)

Let $N \geq 2$ and assume that

$$\Phi_{\text{ext}}(x) = \frac{1}{2} \sum_{j=1}^N \frac{x_j^2}{\lambda_j^2},$$

with $\lambda_j > 0$, $1 \leq j \leq N$. Let us also fix $\mathbf{m} > 0$. Then, there exists a unique $a \in (\mathbb{R}_+^*)^N$ such that

$$(18) \quad |\mathcal{K}_a| \left(\sum_{j=1}^N \frac{1}{\lambda_j^2} \right) = \mathbf{m} \quad \text{and} \quad \frac{\mathbf{m}}{2|B(0,1)| \sum_{k=1}^N \lambda_k^{-2}} \mathcal{Z}(a_1^2, \dots, a_N^2) = \left(\frac{1}{\lambda_j^2} \right)_{1 \leq j \leq N}.$$

Therefore, there exists a constant κ , depending only on the λ_j 's, N and \mathbf{m} such that the function

$$\Phi_e = \Phi_{\text{ext}} + \left(\sum_{j=1}^N \frac{1}{\lambda_j^2} \right) \Gamma \star \mathbf{1}_{\mathcal{K}_a} + \kappa$$

is convex and satisfies

$$(19) \quad -\Delta \Phi_e = \left(\sum_{j=1}^N \frac{1}{\lambda_j^2} \right) \mathbf{1}_{\mathbb{R}^N \setminus \mathcal{K}_a}$$

with, furthermore, $\Phi_e = 0$ in \mathcal{K}_a and $\Phi_e > 0$ in $\mathbb{R}^N \setminus \mathcal{K}_a$.

Proof. — We define $\lambda > 0$ such that $\lambda^{-2} = \sum_{j=1}^N \lambda_j^{-2}$ and the constant κ by the formulas

$$4\kappa = \begin{cases} -\lambda^{-2} \left(\prod_{j=1}^N a_j \right) \int_0^{+\infty} \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds & \text{if } N \geq 3, \\ -\lambda^{-2} (a_1 a_2) \ln((a_1 + a_2)^2 / 2) & \text{if } N = 2. \end{cases}$$

The existence (and uniqueness) of a satisfying the conditions (18) then ensures that $\Phi_e = 0$ in \mathcal{K}_a and $-\Delta \Phi_e = \lambda^{-2} \mathbf{1}_{\mathbb{R}^N \setminus \mathcal{K}_a}$. Then, from the formulas in Proposition 2.1, we get, in $\{\sigma_a > 0\}$,

$$(20) \quad \begin{aligned} \frac{4\lambda^2}{\prod_{j=1}^N a_j} \Phi_e(x) &= \Gamma \star \mathbf{1}_{\mathcal{K}_a}(x) + \frac{4\lambda^2}{\prod_{j=1}^N a_j} \Phi_{\text{ext}}(x) + \frac{4\lambda^2}{\prod_{j=1}^N a_j} \kappa \\ &= \int_{\sigma_a(x)}^{+\infty} \left(1 - \sum_{j=1}^N \frac{x_j^2}{a_j^2 + s} \right) \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds \\ &\quad + \sum_{k=1}^N x_k^2 \int_0^{+\infty} \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} \frac{ds}{a_k^2 + s} - \int_0^{+\infty} \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds \\ &= \int_0^{\sigma_a(x)} \left(\sum_{j=1}^N \frac{x_j^2}{a_j^2 + s} - 1 \right) \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds. \end{aligned}$$

The last integral is positive if $\sigma_a(x) > 0$ since, when $0 \leq s < \sigma_a(x)$,

$$\sum_{j=1}^N \frac{x_j^2}{a_j^2 + s} - 1 > \sum_{j=1}^N \frac{x_j^2}{a_j^2 + \sigma_a(x)} - 1 = 0.$$

In order to see that Φ_e is convex, we notice that $\Phi_e \equiv 0$ in $\{\sigma_a \leq 0\} = \mathcal{K}_a$ and that, from (20), we have, when $\sigma_a(x) > 0$, and for any direction $\omega \in \mathbb{S}^{N-1}$,

$$\begin{aligned} \partial_\omega^2 \Phi_e(x) &= \frac{\prod_{j=1}^N a_j}{2\lambda^2} \int_0^{\sigma_a(x)} \left(\sum_{j=1}^N \frac{\omega_j^2}{a_j^2 + s} \right) \left(\prod_{j=1}^N (a_j^2 + s) \right)^{-1/2} ds \\ &\quad + \frac{\prod_{j=1}^N a_j}{\lambda^2} \left(\sum_{j=1}^N \frac{x_j \omega_j}{a_j^2 + \sigma_a(x)} \right)^2 \left(\sum_{j=1}^N \frac{x_j^2}{(a_j^2 + \sigma_a(x))^2} \right) \left(\prod_{j=1}^N (a_j^2 + \sigma_a(x)) \right)^{-1/2}, \end{aligned}$$

since $\sum_{j=1}^N x_j^2 / (a_j^2 + \sigma_a(x)) = 1$, which is indeed > 0 . □

Clearly, the ellipsoid \mathcal{K}_a is not a level set of the external potential Φ_{ext} (except when all the λ_j 's are all equal). It is interesting to study the limiting case of a very asymmetric external potential. For instance in $N = 2$, we consider a trapping potential (12) with a large aspect ratio $A = \lambda_1/\lambda_2 \gg 1$. Direct computations (using Proposition 2.3) lead to

$$a_1 = \sqrt{\frac{m}{\pi}} \frac{\lambda_1}{\sqrt{1 + \lambda_2^2/\lambda_1^2}}; \quad a_2 = \sqrt{\frac{m}{\pi}} \frac{\lambda_2}{\sqrt{1 + \lambda_1^2/\lambda_2^2}}.$$

Hence

$$a_1/a_2 = \lambda_1^2/\lambda_2^2 = A^2.$$

Thus, the aspect ratio of the particles' cloud is much larger than the aspect ratio of the external potential: this is an effect of the strong repulsion, see Figure 2 for a typical picture. A similar phenomenon occurs in higher dimensions. For $N = 3$ with cylindrical symmetry, explicit formulas corresponding to our \mathcal{Z} function are given for instance in [47]. It is easy to check that for a strongly oblate external potential (“pancake shape”), the aspect ratio of the cloud is again of the order of the square of the aspect ratio of the external potential. We can now state the analog of Theorem 1.2 for a general quadratic Φ_{ext} .

THEOREM 2.6. — *Let Φ_{ext} be any quadratic potential (12) to which we associate, by virtue of Corollary 2.5, the ellipsoid \mathcal{K}_a and the potential Φ_e . Let $V^{\text{init}} \in H^s(\overset{\circ}{\mathcal{K}}_a)$ satisfy $\nabla_x \cdot V^{\text{init}} = 0$ in $\overset{\circ}{\mathcal{K}}_a$ and the no-flux condition (6) on $\partial\mathcal{K}_a$. Denote by V the solution on $[0, T]$ to (IE) with the no-flux condition (6) given in Theorem 1.1 and consider $\mathcal{V}^{\text{init}}$ a smooth extension of V in \mathbb{R}^N satisfying the following conditions, where $R > 0$ is such that $\mathcal{K}_a \subset B(0, R)$,*

$$\mathcal{V}|_{\mathcal{K}_a} = V, \quad \mathcal{V}|_{\mathbb{R}^N \setminus B(0, 2R)} = 0, \quad \mathcal{V}(t, x) \cdot \nu(x)|_{\partial\mathcal{K}_a} = 0.$$

Let $f_\varepsilon^{\text{init}} : \mathbb{R}^N \times \mathbb{R}^N \rightarrow [0, \infty)$ be a sequence of integrable functions that satisfy (11). Then, the associated solution f_ε of the Vlasov–Poisson equation (V)–(P) satisfies, as $\varepsilon \rightarrow 0$,

- (i) ρ_ε converges to $n_e = (\sum_{j=1}^N \lambda_j^{-2}) \mathbf{1}_{\mathcal{K}_a}$ in $C^0(0, T; \mathcal{M}^1(\mathbb{R}^N)\text{-weak-}\star)$;
- (ii) $\mathcal{H}_{\mathcal{V}, \varepsilon}$ converges to 0 uniformly on $[0, T]$;

(iii) J_ε converges to J in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$ weakly-*, the limit J lies in $L^\infty(0, T; L^2(\mathbb{R}^N))$ and satisfies $J|_{[0, T] \times \mathcal{K}_a} = V$, $\nabla_x \cdot J = 0$ and $J \cdot \nu(x)|_{\partial \mathcal{K}_a} = 0$.

The existence of a smooth extension \mathcal{V} of V on \mathbb{R}^N satisfying the above mentioned constraints follows from [33, Chap. I: Th. 2.1 p. 17 & Th. 8.1 p. 42]. See Lemma B.1 in the appendix for a divergence-free extension.

2.2. THE CASE OF A GENERAL POTENTIAL. — We wish now to extend the above results to a general confining potential. When Φ_{ext} is not quadratic, the equilibrium density n_e cannot be expected to be constant on its support. In turn, the limiting equation will be more complicated than the Incompressible Euler system. Besides, the determination of the domain $\mathcal{K} = \{\Phi_e = 0\}$ is a non trivial issue, and its geometry might be quite involved [43]. In the following we write $\Omega = \mathcal{K}$ for the interior of \mathcal{K} .

The pair (n_e, Ω) should be thought of through energetic consideration. As it will be detailed below, the total energy of the system (V)-(P) is

$$\iint \frac{|v|^2}{2} f_\varepsilon \, dv \, dx + \frac{1}{2\varepsilon} \int \Phi_\varepsilon \rho_\varepsilon \, dx + \frac{1}{\varepsilon} \int \Phi_{\text{ext}} \rho_\varepsilon \, dx.$$

It is natural to investigate solutions whose energy does not diverge when ε tends to 0. Hence we are interested in configurations close to the ground state n_e defined by the variational problem where only the electrostatic part of the energy is involved: namely, we wish to minimize

$$\mathcal{E}[\rho] := \int \Phi_{\text{ext}}(x) \, d\rho(x) + \frac{1}{2} \iint \Gamma(x - y) \, d\rho(y) \, d\rho(x),$$

for a fixed $\mathbf{m} > 0$ over the convex subset $\mathcal{M}_{\text{ext}}^+(\mathbf{m})$ made of nonnegative Borel measures ρ of total mass $\mathbf{m} > 0$ such that $\int \Phi_{\text{ext}} \, d\rho$ is finite. This problem, which is often referred to as the generalized Gauss variational problem, is quite classical and the basis of the theory dates back to [18]. We refer the reader to [42, Chap. 1] for the case $N = 2$, and to [8, Th. 1.2] when $N \geq 3$ for the existence of a minimizer under suitable assumptions on Φ_{ext} . In what follows, we shall assume that Φ_{ext} fulfills the following requirements:

(h1) $\Phi_{\text{ext}} : \mathbb{R}^N \rightarrow \mathbb{R}_+$ is continuous, nonnegative and satisfies $\Phi_{\text{ext}}(x) \rightarrow +\infty$ as $|x| \rightarrow +\infty$,

(h2) If $N = 2$ or $N = 1$, we have $\lim_{|x| \rightarrow +\infty} (\Phi_{\text{ext}} + \mathbf{m}\Gamma)(x) = +\infty$.

The following statement collects from [8, 42, 43] the results we shall need.

THEOREM 2.7. — *We assume that the potential Φ_{ext} satisfies the hypotheses (h1) and (h2).*

(i) *The functional \mathcal{E} is strictly convex on $\mathcal{M}_{\text{ext}}^+(\mathbf{m})$.*

(ii) *The problem*

$$(21) \quad \inf \{ \mathcal{E}[\rho] ; \rho \in \mathcal{M}_{\text{ext}}^+(\mathbf{m}) \}$$

has a unique minimizer n_e which has a compact support of positive capacity. Moreover, there exists a constant C_* such that

$$(22) \quad \begin{cases} \Gamma \star n_e + \Phi_{\text{ext}} \geq C_* & \text{quasi everywhere,} \\ \Gamma \star n_e + \Phi_{\text{ext}} = C_* & \text{quasi everywhere on } \text{Supp}(n_e). \end{cases}$$

(iii) Conversely, assume that $\rho_0 \in \mathcal{M}_{\text{ext}}^+(\mathfrak{m})$ and C_0 are such that

$$\begin{cases} \Gamma \star \rho_0 + \Phi_{\text{ext}} \geq C_0 & \text{quasi everywhere,} \\ \Gamma \star \rho_0 + \Phi_{\text{ext}} = C_0 & \text{quasi everywhere on } \text{Supp}(\rho_0). \end{cases}$$

Then, ρ_0 is the minimizer for (21): $\rho_0 = n_e$.

We then define the potential $\Phi_e := \Gamma \star n_e + \Phi_{\text{ext}} - C_*$. The constant C_* in (22) is called the modified Robin constant and *quasi everywhere* (q. e.) means up a set of zero capacity (which is a bit stronger than to be Lebesgue-negligible); see [43, Def. 2.11]. If $N = 1$, (22) holds pointwise.

Proof. — The statements for $N \geq 3$ can be found in [8, Th.1.2]. When $N = 2$, we refer the reader to [42, Th. 1.3 for (ii) and Th. 3.3 for (iii)]. If $N = 2$, the strict convexity (i) is not made explicit in [42]. Thus we give proofs of (i) for $N = 2$, and (i)–(iii) for $N = 1$.

The argument for (i) is that if $\rho_0, \rho_1 \in \mathcal{M}_{\text{ext}}^+(\mathfrak{m})$ and $\theta \in (0, 1)$, then

$$\begin{aligned} & \mathcal{E}[(1 - \theta)\rho_0 + \theta\rho_1] - (1 - \theta)\mathcal{E}[\rho_0] - \theta\mathcal{E}[\rho_1] \\ &= \frac{1}{2} \iint \Gamma(x - y) \, d[(1 - \theta)\rho_0 + \theta\rho_1](y) \, d[(1 - \theta)\rho_0 + \theta\rho_1](x) \\ & \quad - \frac{1}{2}(1 - \theta) \iint \Gamma(x - y) \, d\rho_0(y) \, d\rho_0(x) - \frac{1}{2}\theta \iint \Gamma(x - y) \, d\rho_1(y) \, d\rho_1(x) \\ &= -\frac{1}{2}\theta(1 - \theta) \iint \Gamma(x - y) \, d[\rho_0 - \rho_1](y) \, d[\rho_0 - \rho_1](x). \end{aligned}$$

Unless $\rho_0 = \rho_1$, the last integral is shown to be positive if $N \geq 3$ in [8, Lem. 3.1]. The case $N = 2$ is dealt with in [42, Lem. 1.8], under the restriction that $\rho_0 - \rho_1$ has compact support. Actually, the method used in [8], which consists in writing $\Gamma(x)$ as an integral of Gaussians $e^{-|x|^2/2t}$, can be extended to the case $N = 2$ as we check now.

The starting point is the equality (see [2, Eq. (12)])

$$\ln \frac{1}{r} = \int_0^{+\infty} \frac{1}{2t} \left(e^{-r^2/2t} - e^{-1/2t} \right) dt.$$

Therefore, denoting $\rho := \rho_0 - \rho_1$ and $r := |x - y|$ and using the dominated convergence theorem (on each of the sets $\{|x - y| < 1\}$ and $\{|x - y| \geq 1\}$), we obtain

$$\begin{aligned} \iint \ln \frac{1}{|x - y|} d\rho(y) d\rho(x) &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{2t} \iint (e^{-r^2/2t} - e^{-1/2t}) d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{2t} \iint e^{-r^2/2t} d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{4\pi} \iiint e^{-t|\xi|^2/2 - i\xi \cdot (x-y)} d\xi d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \int \frac{1}{4\pi} e^{-t|\xi|^2/2} |\widehat{\rho}(\xi)|^2 d\xi dt = \int \frac{1}{2\pi|\xi|^2} |\widehat{\rho}(\xi)|^2 d\xi, \end{aligned}$$

where, for the second equality, we use $\int d\rho = 0$, and for the third one, we write $e^{-r^2/2t}$ as the Fourier transform of a two dimensional Gaussian. This clearly shows that $\iint \Gamma(x - y) d\rho(y) d\rho(x)$ is positive unless $\rho = 0$, ensuring the strict convexity of \mathcal{E} on $\mathcal{M}_{\text{ext}}^+(\mathfrak{m})$. When $N = 1$, we argue in a similar way by observing that

$$-r = \int_0^{+\infty} \frac{1}{\sqrt{2\pi t}} (e^{-r^2/2t} - 1) dt.$$

Indeed,

$$e^{-r^2/2t} - 1 = \int_0^r \partial_u (e^{-u^2/2t}) du = - \int_0^r (u/t) e^{-u^2/2t} du,$$

thus

$$\begin{aligned} - \int_0^{+\infty} \frac{1}{\sqrt{t}} (e^{-r^2/2t} - 1) dt &= \int_0^{+\infty} \frac{1}{\sqrt{t}} \int_0^r (u/t) e^{-u^2/2t} du dt \\ &= \int_0^r \int_0^{+\infty} \frac{u}{t^{3/2}} e^{-u^2/2t} dt du \\ &= \int_0^r \int_0^{+\infty} 2\sqrt{2} e^{-\tau^2} d\tau du = \int_0^r \sqrt{2\pi} du = r \sqrt{2\pi}, \end{aligned}$$

where we have used the change of variable $\tau = u/\sqrt{2t}$. Owing to this relation, we can follow the same lines as above:

$$\begin{aligned} -\frac{1}{2} \iint |x - y| d\rho(y) d\rho(x) &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{2\sqrt{2\pi t}} \iint (e^{-r^2/2t} - 1) d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{2\sqrt{2\pi t}} \iint e^{-r^2/2t} d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \frac{1}{4\pi} \iiint e^{-t\xi^2/2 - i\xi \cdot (x-y)} d\xi d\rho(y) d\rho(x) dt \\ &= \lim_{T \rightarrow +\infty} \int_{1/T}^T \int \frac{1}{4\pi} e^{-t\xi^2/2} |\widehat{\rho}(\xi)|^2 d\xi dt = \int \frac{1}{2\pi|\xi|^2} |\widehat{\rho}(\xi)|^2 d\xi. \end{aligned}$$

It only remains to prove (ii)–(iii) for $N = 1$. This is tackled in [43], but with $\Gamma(x) = -\ln|x|$. The very same arguments apply to the case $\Gamma(x) = -|x|/2$. \square

REMARK 2.8. — To motivate the above computation, one can remark that for $N \geq 1$ and under the condition $\int d\rho = 0$, we have, at least formally,

$$\iint \Gamma(x - y) d\rho(y) d\rho(x) = \int |\nabla \Delta^{-1} \rho|^2 dx = (2\pi)^{-N} \int \frac{1}{|\xi|^2} |\widehat{\rho}(\xi)|^2 d\xi.$$

The minimization of the functional \mathcal{E} is connected to an *obstacle problem*. This connection is explained in details in [43, §2.5].

PROPOSITION 2.9. — *If n_e is the minimizer of Theorem 2.7, then $h = \Gamma \star n_e$ is the unique solution to the obstacle problem:*

To find $\phi \in H_{\text{loc}}^1(\mathbb{R}^N)$ such that

$$\int \nabla \phi \cdot \nabla (g - \phi) dx \geq 0,$$

holds for any $g \in H_{\text{loc}}^1(\mathbb{R}^N)$, with $g - \phi$ compactly supported and $\phi \geq \psi$ q. e.

where $\psi(x) := C_* - \Phi_{\text{ext}}(x)$. Furthermore, assume $\Phi_{\text{ext}} \in C^{1,1}(\mathbb{R}^N)$. Then, $\Gamma \star n_e \in C^{1,1}(\mathbb{R}^N)$.

The second part of the statement, which improves the natural regularity of $\Gamma \star n_e$, is due to [6], [15], [26] with an analysis that relies on PDEs tools for the obstacle problem, instead of dealing directly with the minimization problem (21). We then define the coincidence set

$$\mathcal{K} := \{\Phi_e = 0\} = \{\Gamma \star n_e = C_* - \Phi_{\text{ext}}\}$$

and claim that \mathcal{K} is compact. Indeed, as $|x| \rightarrow +\infty$, we have $-\Gamma \star n_e(x) \sim -\mathbf{m}\Gamma(x)$ and $\Phi_{\text{ext}}(x) + \mathbf{m}\Gamma(x) \gg 1$ whatever is the dimension N by (h1)–(h2), thus $\mathcal{K} = \{\Gamma \star n_e = C_* - \Phi_{\text{ext}}\}$ is bounded. The regularity of $\Gamma \star n_e$ implies that \mathcal{K} is closed as the level set of a continuous function.

Moreover, by (22), the set $\text{Supp}(n_e) \setminus \mathcal{K}$ has zero capacity. We give some examples in section 2.3 below where $\text{Supp}(n_e) \subsetneq \mathcal{K}$, due to the presence of points or regions where $\Delta \Phi_{\text{ext}}$ vanishes. These points are precisely defined in [26, §3.6] and called ‘shallow points’ and it is shown in this paper (see Proposition 3.12 there) that it is possible to pass from $\text{Supp}(n_e)$ to \mathcal{K} by simply adding these ‘shallow points’. This fact is illustrated in section 2.3 below.

For a general potential Φ_{ext} , the variational viewpoint and the theory of the obstacle problem provide a definition for the equilibrium distribution n_e , the domain \mathcal{K} (which is not always the support of n_e) and the potential Φ_e . The regularity of Φ_{ext} is not “transferred” to the solution Φ_e or $\Gamma \star n_e$ beyond $C^{1,1}$ regularity (see [15], [6]) since the Laplacian of these functions is discontinuous. In addition, the topology of \mathcal{K} is difficult to analyze in general: \mathcal{K} may have empty interior or may exhibit cusps. Hence, these regularity issues for both \mathcal{K} and n_e need to be discussed individually. Let us then list the properties, which very likely are far from optimal, that we need to deal with the asymptotic regime: there exists $s > 1 + N/2$ such that

(H1) \mathcal{K} has a non empty interior Ω and $\partial\Omega$ is of class C^1 .

(H2) $\Phi_{\text{ext}} \in C^{s+3}(\mathbb{R}^N)$, $\Delta \Phi_{\text{ext}}$ is bounded away from zero on \mathcal{K} .

The C^1 regularity assumption (H1) on $\partial\Omega$ excludes the presence of cusps in \mathcal{K} . Let us point out a regularity result derived from the obstacle problem theory.

PROPOSITION 2.10. — *Let Φ_{ext} be a potential satisfying (h1) and (h2) and consider the minimizer n_e of (21) given by Theorem 2.7 and let $\mathcal{K} := \{\Phi_e = 0\}$. Assume that (H1) and (H2) are satisfied.*

- (i) [29] *Then, $\bar{\Omega} = \mathcal{K}$ and the boundary $\partial\Omega$ is C^{s+1} .*
- (ii) *We have $n_e = \mathbf{1}_\Omega(\Delta\Phi_{\text{ext}})$ as measures.*

REMARK 2.11. — A consequence of the $C^{1,1}$ regularity of $\Gamma \star n_e$ is that, in (22), we may replace ‘quasi everywhere’ by ‘everywhere’ since all the functions involved are continuous. In particular, note that the strengthened assumption (H2) then implies that n_e is $C^{s+1}(\bar{\Omega})$ and bounded from below on \mathcal{K} .

REMARK 2.12. — Under the low regularity assumption $\Phi_{\text{ext}} \in C^1(\mathbb{R}^N)$ and when $N \geq 2$, it follows from [6, Th. 2] that $\Gamma \star n_e \in C^1$. This prevents the singular part of the measure n_e from being a Dirac mass or a finite sum of Dirac masses, since the fundamental solution Γ is unbounded (if $N \geq 2$) near the origin. This however may happen in dimension 1 (see the examples in section 2.3 below) and in these cases, the relation $n_e = \mathbf{1}_\Omega(\Delta\Phi_{\text{ext}})$ (as measures) might not be true.

Proof. — For the first point, notice that $\bar{\Omega} = \mathcal{K}$ by (H1). In addition, we also have from [29, Th. 1] or [16, Chap. 2, Th. 1.1], since we assume that $\Delta\Phi_{\text{ext}} \in C_{\text{loc}}^{s+1}(\mathbb{R}^N)$ and does not vanish on \mathcal{K} , that the boundary $\partial\Omega$ is automatically of class C^{s+1} (and even $C^{s+1+\beta}$ for any $\beta \in (0, 1)$).

For the second statement, we first invoke the regularity result of [15] (see also [6], [16]) saying that since $\Phi_{\text{ext}} \in C^{s+3}(\mathbb{R}^N) \subset C^{1,1}(\mathbb{R}^N)$, then $\Gamma \star n_e$ belongs to $C^{1,1}(\mathbb{R}^N)$. Consequently, in the distributional sense, the compactly supported measure $n_e = -\Delta(\Gamma \star n_e)$ belongs to $L^\infty(\mathbb{R}^N)$. Since $\Gamma \star n_e = C_* - \Phi_{\text{ext}}$ in Ω (q. e., hence everywhere by continuity of the functions), we infer that $n_e = \Delta\Phi_{\text{ext}}$ in Ω . If we make the assumption (H1), $\partial\Omega$ is of class C^1 and we then deduce immediately that $n_e = (\Delta\Phi_{\text{ext}})\mathbf{1}_\Omega$ as a measure. When assumption (H1) is not satisfied, then, as noticed in [26, Th. 3.10], it follows from [30, Chap. 2, Lem. A.4] that $n_e = \Delta\Phi_{\text{ext}}$ holds almost everywhere in \mathcal{K} , which concludes, even in this case. \square

REMARK 2.13. — The arguments in [16] and [29] can be used to establish that if Φ_{ext} is analytic, then $\partial\Omega$ is analytic too, when (H1) and (H2) hold.

With these assumptions (H1) and (H2), we can establish the following statement, where we point out that the limit problem is the Lake Equation (LE) instead of the mere incompressible Euler system, since now the equilibrium distribution n_e is inhomogeneous. We obviously need a smooth enough solution to the Lake Equation (LE): we may refer to the works [32], [39] (when the domain Ω is simply connected) and [31] (without simple connectedness assumption on the domain Ω), which rely on a vorticity formulation *à la* Yudovitch and are then restricted to the dimension $N = 2$.

We provide in the appendix (see Theorem A.1) a well-posedness result analogous to Theorem 1.1 valid in any dimension and without simple connectedness assumption on Ω .

THEOREM 2.14. — *Let Φ_{ext} be a potential satisfying (h1) and (h2) and consider n_e the minimizer of (21) given by Theorem 2.7 and let $\mathcal{K} := \{\Phi_e = 0\}$. Assume in addition that (H1) and (H2) are satisfied. Let $V^{\text{init}} \in H^s(\Omega)$ satisfy $\nabla_x \cdot (n_e V^{\text{init}}) = 0$ in Ω and the no-flux condition (6). Denote by V the solution on $[0, T]$ to the Lake Equation (LE), with the no-flux condition (6) and initial condition V^{init} , given in Theorem A.1 and consider $\mathcal{V}^{\text{init}}$ a smooth extension of V^{init} satisfying the following conditions, where $R > 0$ is such that $\Omega \subset B(0, R)$,*

$$\mathcal{V}|_{\Omega} = V, \quad \mathcal{V}|_{\mathbb{R}^N \setminus B(0, 2R)} = 0, \quad \mathcal{V}(t, x) \cdot \nu(x)|_{\partial\Omega} = 0.$$

Let $f_{\varepsilon}^{\text{init}} : \mathbb{R}^N \times \mathbb{R}^N \rightarrow [0, \infty)$ be a sequence of integrable functions that satisfy (11). Then, the associated solution f_{ε} of the Vlasov–Poisson equation (V)–(P) satisfies, as $\varepsilon \rightarrow 0$,

- (i) ρ_{ε} converges to n_e in $C^0(0, T; \mathcal{M}^1(\mathbb{R}^N))$ -weak- \star ;
- (ii) $\mathcal{H}_{\mathcal{V}, \varepsilon}$ converges to 0 uniformly on $[0, T]$;
- (iii) J_{ε} converges to J in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$ weakly- \star , the limit J lies in $L^{\infty}(0, T; L^2(\mathbb{R}^N))$ and satisfies $J|_{[0, T] \times \Omega} = V$, $\nabla_x \cdot J = 0$ and $J \cdot \nu(x)|_{\partial\Omega} = 0$.

The existence of a smooth extension \mathcal{V} of V follows from [33, Chap. I: Th. 2.1 p. 17 & Th. 8.1 p. 42].

For convex potentials Φ_{ext} , the only situation where we have been able to check the hypotheses (H1) and (H2) (except the quadratic potentials for which $\Delta\Phi_{\text{ext}}$ is constant) is the case of the space dimension $N = 1$ (see Proposition 2.15) and the case of a radial potential (see Proposition 2.16 below).

2.3. ABOUT HYPOTHESIS (H1) FOR CONVEX POTENTIALS Φ_{ext} . — For the problem we have in mind, it is natural to assume that the confining potential Φ_{ext} is smooth and convex. In this case, one may think that the coincidence set \mathcal{K} or $\text{Supp}(n_e)$ is convex. We have not been able to find such a result in the literature for a general convex, coercive and smooth enough confining potential Φ_{ext} . Actually, the obstacle problem is, in most cases, set on a bounded convex domain G with suitable boundary conditions instead of the whole space \mathbb{R}^N .

For the obstacle problem in bounded convex domains G , we can find a convexity result for the coincidence set \mathcal{K} in [17, Th. 6.1] in the specific assumptions that $\Delta\Phi_{\text{ext}}$ is constant and with the boundary condition $\Gamma \star n_e = 1 + \psi = 1 + C_* - \Phi_{\text{ext}}$ on ∂G . Just after [17, Th. 6.1], an example is given (in a bounded convex domain G) showing that the assumption Φ_{ext} smooth and strictly convex (and $\Gamma \star n_e > \psi$ on ∂G) is not sufficient to guarantee that \mathcal{K} is convex. Roughly speaking, $\Delta\Phi_{\text{ext}}$ is constant for quadratic potentials.

Turning back to the obstacle problem in the whole space \mathbb{R}^N , the only convexity result we are aware of is [6, Cor. 7], which corresponds, again, to the case where $\Delta\Phi_{\text{ext}}$

is constant. Extending this result to space depending functions $\Delta\Phi_{\text{ext}}$ is a delicate issue (see however [16, Chap. 2, §3], which is not sufficient for our situation).

In the one dimensional case and for a convex potential Φ_{ext} , there is a simple characterization of \mathcal{K} , as explicited in the following Proposition.

PROPOSITION 2.15 (The one dimensional case with a convex potential)

Assume that $N = 1$ and that $\Phi_{\text{ext}} : \mathbb{R} \rightarrow \mathbb{R}$ is of class C^1 , piecewise C^2 , nonnegative, convex (i.e., Φ'_{ext} is nondecreasing) and that $\Phi_{\text{ext}}(x) - \mathfrak{m}|x|/2 \gg 1$ for $|x| \gg 1$ (so that (h1) and (h2) are satisfied). We denote by $\partial\Phi'_{\text{ext}}$ the piecewise continuous function associated with the second order derivative of Φ_{ext} . Then, the minimizer n_e for (21) is given by

$$(23) \quad n_e = (\partial\Phi'_{\text{ext}})|_{]a_-, a_+[},$$

where a_+ and a_- are defined by the equations

$$(24) \quad \mathfrak{m}/2 = \Phi'_{\text{ext}}(a_+) \quad \text{and} \quad -\mathfrak{m}/2 = \Phi'_{\text{ext}}(a_-).$$

Furthermore, $\text{Supp}(n_e) = \text{Supp}(\partial\Phi'_{\text{ext}}) \cap [a_-, a_+]$ and $\{\Phi_e = 0\} = [a_-, a_+]$. In addition, the potential Φ_e is convex.

Proof. — As a first observation, notice that Φ'_{ext} is continuous, nondecreasing and admits limits: $\lim_{x \rightarrow +\infty} \Phi'_{\text{ext}}(x) \geq \mathfrak{m}/2$, and $\lim_{x \rightarrow -\infty} \Phi'_{\text{ext}}(x) \leq -\mathfrak{m}/2$. Assume $\lim_{x \rightarrow +\infty} \Phi'_{\text{ext}}(x) = \mathfrak{m}/2$. The convexity of Φ_{ext} implies $\Phi_{\text{ext}}(x) - \Phi_{\text{ext}}(y) \leq \frac{\mathfrak{m}}{2}(x - y)$ for $x > y$. Letting x go to $+\infty$ contradicts the assumption on the asymptotic behavior of Φ_{ext} . We infer that $\lim_{x \rightarrow +\infty} \Phi'_{\text{ext}}(x) > \mathfrak{m}/2$. A similar reasoning shows that $\lim_{x \rightarrow -\infty} \Phi'_{\text{ext}}(x) < -\mathfrak{m}/2$. We conclude that (24) has at least one (possibly non unique) solution.

Since Φ'_{ext} is nondecreasing, and if $b_+ > a_+$ also solves $\mathfrak{m}/2 = \Phi'_{\text{ext}}(b_+)$, this implies that, on $[a_+, b_+]$, $\Phi'_{\text{ext}} \equiv \mathfrak{m}/2$, thus $\partial\Phi'_{\text{ext}} \equiv 0$ and this does not change n_e .

Let us use the characterization (iii) in Theorem 2.7 and look for the measure n_e under the form $n_e = (\partial\Phi'_{\text{ext}})|_{]a_-, a_+[}$, which is piecewise continuous. This function n_e satisfies the mass constraint if and only if

$$(25) \quad \mathfrak{m} = \int_{a_-}^{a_+} \partial\Phi'_{\text{ext}} dx = \Phi'_{\text{ext}}(a_+) - \Phi'_{\text{ext}}(a_-).$$

Now, let us compute $\Gamma \star n_e + \Phi_{\text{ext}}$ in $[a_-, a_+]$ and investigate under which condition this function is constant (in $[a_-, a_+]$). Elementary computations give, for $a_- \leq x \leq a_+$:

$$\begin{aligned} \Gamma \star n_e(x) &= -\frac{1}{2} \int_{a_-}^{a_+} |y - x| (\partial\Phi'_{\text{ext}})(y) dy \\ &= -\frac{1}{2} \Phi'_{\text{ext}}(a_+)(a_+ - x) + \frac{1}{2} \Phi'_{\text{ext}}(a_-)(x - a_-) + \frac{1}{2} \int_{a_-}^{a_+} \text{sgn}(y - x) \Phi'_{\text{ext}}(y) dy \\ &= -\frac{1}{2} \Phi'_{\text{ext}}(a_+)(a_+ - x) + \frac{1}{2} \Phi'_{\text{ext}}(a_-)(x - a_-) + \frac{1}{2} \Phi_{\text{ext}}(a_+) \\ &\quad + \frac{1}{2} \Phi_{\text{ext}}(a_-) - \Phi_{\text{ext}}(x). \end{aligned}$$

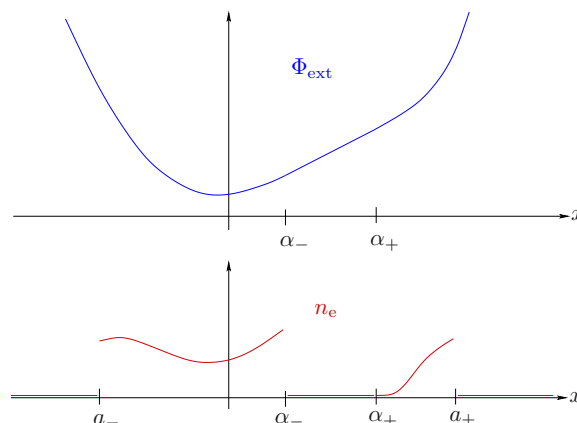


FIGURE 3. The potential Φ_{ext} and the corresponding measure n_e for Example 2

As a consequence, $\Gamma \star n_e + \Phi_{\text{ext}}$ is constant in $[a_-, a_+]$ if and only if $\Phi'_{\text{ext}}(a_+) + \Phi'_{\text{ext}}(a_-) = 0$. Combining this with the mass constraint $\Phi'_{\text{ext}}(a_+) - \Phi'_{\text{ext}}(a_-) = \mathbf{m}$ yields the relation (24). It then follows that, on $[a_-, a_+]$,

$$\begin{aligned} \Gamma \star n_e + \Phi_{\text{ext}} &= C_* := \frac{1}{2} (\Phi_{\text{ext}}(a_+) + \Phi_{\text{ext}}(a_-) - a_+ \Phi'_{\text{ext}}(a_+) - a_- \Phi'_{\text{ext}}(a_-)) \\ &= \frac{1}{2} (\Phi_{\text{ext}}(a_+) + \Phi_{\text{ext}}(a_-)) - \frac{\mathbf{m}}{4} (a_+ - a_-). \end{aligned}$$

It remains to check that $\Gamma \star n_e + \Phi_{\text{ext}} \geq C_*$ in \mathbb{R} . To see this, note that $\Phi_e := \Gamma \star n_e + \Phi_{\text{ext}} - C_*$ is convex since its (distributional) second order derivative is equal to the piecewise continuous function $\partial_x \Phi'_{\text{ext}} \mathbf{1}_{\mathbb{R} \setminus [a_-, a_+]}$, and $\Phi_e \equiv 0$ on $[a_-, a_+]$, hence is ≥ 0 everywhere. This finishes the proof. \square

Let us give some examples illustrating Proposition 2.15.

EXAMPLE 1 (1D). — If Φ_{ext} is of class C^2 and Φ''_{ext} is positive on \mathbb{R} , then $n_e(x) = \Phi''_{\text{ext}}(x) \mathbf{1}_{[a_-, a_+]}(x)$ and is absolutely continuous with respect to the Lebesgue measure. We then have $\text{Supp}(n_e) = [a_-, a_+]$.

EXAMPLE 2 (1D). — The potential Φ_{ext} is C^1 , piecewise C^2 , but is affine on the interval $[\alpha_-, \alpha_+]$ (hence it is not strictly convex), where its slope belongs to $] -\mathbf{m}/2, +\mathbf{m}/2[$ (see Figure 3). In addition, the second order derivative Φ''_{ext} is discontinuous at α_- and continuous at α_+ and $\partial \Phi'_{\text{ext}}$ is positive except on $[\alpha_-, \alpha_+]$. In this case, we may still define a_{\pm} as the unique solutions to $\Phi'_{\text{ext}}(a_{\pm}) = \pm \mathbf{m}/2$, and we have $\text{Supp}(n_e) = [a_-, \alpha_-] \cup [\alpha_+, a_+] \subsetneq [a_-, a_+] = \{\Phi_e = 0\}$ and this is then a disconnected set. If the slope in the region $[\alpha_-, \alpha_+]$ where Φ_{ext} is affine does not belong to $] -\mathbf{m}/2, +\mathbf{m}/2[$, then the support of n_e is an interval as in Example 1.

Example 1 fits the hypotheses of Theorem 2.14, but not Example 2 since n_e is not bounded away from zero (near α_+). In particular, for Example 2, we have to

face new difficulties in solving the Cauchy problem (see Theorem A.1) for the Lake Equation (LE). If in the one dimensional situation one can easily check that the support of n_e (instead of \mathcal{X}) is smooth, in a similar higher dimensional case, the regularity of $\text{Supp}(n_e)$ is certainly not easy to analyze since we can not rely on the results in [29, Th. 1] or [16, Chap. 2, Th. 1.1]. All these issues motivate hypothesis (H2).

Let us give now examples which do not fit the regularity hypotheses required in Proposition 2.15. These expressions are justified through the characterization (iii) in Theorem 2.7 and simple computation of $\Gamma \star n_e$.

EXAMPLE 3 (1D). — Take the potential $\Phi_{\text{ext}}(x) = |x|$. Then, hypothesis (h2) exactly means $\mathbf{m} < 1$. In that case, we have $n_e = \mathbf{m}\delta_0$ and $\text{Supp}(n_e) = \{0\} = \{\Phi_e = 0\}$.

EXAMPLE 4 (1D). — Take two reals $a < b$ and a convex potential Φ_{ext} which is affine on $] - \infty, a]$, on $[a, b]$ and on $[b, +\infty[$. Assume also that (h2) is satisfied, that is,

$$\mathbf{m} < \min(\Phi'_{\text{ext}}(+\infty), -\Phi'_{\text{ext}}(-\infty)).$$

Then,

$$n_e = \min\left(\frac{1}{2}\left(\mathbf{m} + \frac{\Phi_{\text{ext}}(b) - \Phi_{\text{ext}}(a)}{b - a}\right)_+, \mathbf{m}\right)\delta_a + \min\left(\frac{1}{2}\left(\mathbf{m} - \frac{\Phi_{\text{ext}}(b) - \Phi_{\text{ext}}(a)}{b - a}\right)_+, \mathbf{m}\right)\delta_b.$$

As a consequence:

- if $-\mathbf{m} < \frac{\Phi_{\text{ext}}(b) - \Phi_{\text{ext}}(a)}{b - a} < \mathbf{m}$, then $\text{Supp}(n_e) = \{a, b\}$ and $\{\Phi_e = 0\} = [a, b]$;
- if $\frac{\Phi_{\text{ext}}(b) - \Phi_{\text{ext}}(a)}{b - a} \leq -\mathbf{m}$, then $n_e = \mathbf{m}\delta_b$ and $\text{Supp}(n_e) = \{b\} = \{\Phi_e = 0\}$;
- if $\frac{\Phi_{\text{ext}}(b) - \Phi_{\text{ext}}(a)}{b - a} \geq \mathbf{m}$, then $n_e = \mathbf{m}\delta_a$ and $\text{Supp}(n_e) = \{a\} = \{\Phi_e = 0\}$.

EXAMPLE 5 (1D). — Consider the potential $\Phi_{\text{ext}}(x) = |x| + x^2/2 + \max(x - 1, 0)$:

- if $\mathbf{m} \leq 2$, then $n_e = \mathbf{m}\delta_0$, $\Gamma \star n_e(x) = -\mathbf{m}|x|/2$ and

$$\text{Supp}(n_e) = \{0\} = \{\Phi_e = 0\};$$

- if $2 \leq \mathbf{m} \leq 4$, then $n_e = 2\delta_0 + \mathbf{1}_{[-\mathbf{m}/2+1, \mathbf{m}/2-1]}$ and

$$\text{Supp}(n_e) = [-\mathbf{m}/2 + 1, \mathbf{m}/2 - 1] = \{\Phi_e = 0\};$$

- if $4 \leq \mathbf{m} \leq 6$, then $n_e = 2\delta_0 + (\mathbf{m}/2 - 2)\delta_1 + \mathbf{1}_{[-\mathbf{m}/2+1, 1]}$ and

$$\text{Supp}(n_e) = [-\mathbf{m}/2 + 1, 1] = \{\Phi_e = 0\};$$

- if $\mathbf{m} \geq 6$, then $n_e = 2\delta_0 + \delta_1 + \mathbf{1}_{[-\mathbf{m}/2+1, \mathbf{m}/2-2]}$ and

$$\text{Supp}(n_e) = [-\mathbf{m}/2 + 1, \mathbf{m}/2 - 2] = \{\Phi_e = 0\}.$$

Examples 3, 4 and 5 show that the single convexity hypothesis on Φ_{ext} does not guarantee that n_e is a restriction of the nonnegative measure $\partial_x^2 \Phi_{\text{ext}}$ (in the distributional sense). It appears in these examples that n_e is nondecreasing with respect to the mass \mathbf{m} , and thus that we always have $n_e \leq \partial_x^2 \Phi_{\text{ext}}$ in the distributional sense. It

is an open problem to determine whether this holds true in higher dimensions. Here again, these issues motivate the regularity assumptions on Φ_{ext} in (H2).

The other situation where we may verify hypothesis (H1) is the radial case (see [8, Cor.1.4] for a related result in dimension $N \geq 3$ for C^2 potentials Φ_{ext}). Let $\varphi_{\text{ext}} : \mathbb{R}_+ \rightarrow \mathbb{R}_+$ be a nondecreasing function of class C^1 and piecewise C^2 . Consider now the potential $\Phi_{\text{ext}} : \mathbb{R}^N \rightarrow \mathbb{R}$ given by $\Phi_{\text{ext}}(x) = \varphi_{\text{ext}}(|x|)$. It is then clear that φ_{ext} is convex if and only if Φ_{ext} is convex.

PROPOSITION 2.16 (The radial case with a convex potential)

Assume that $N \geq 2$ and that $\Phi_{\text{ext}} : \mathbb{R}^N \rightarrow \mathbb{R}$ is as above. Then, the minimizer n_e for (21) is given by

$$(26) \quad n_e(x) = \mathbf{1}_{B(0,R)}(x) \Delta \Phi_{\text{ext}}(x),$$

where R is defined by the equation

$$(27) \quad \mathbf{m} = \int_{B(0,R)} \Delta \Phi_{\text{ext}}(x) \, dx \quad \text{or, equivalently,} \quad N|B(0,1)|R^{N-1} \varphi'_{\text{ext}}(R) = \mathbf{m}.$$

Furthermore, $\text{Supp}(n_e) = \overline{B}(0, R) \setminus B(0, R_{\min})$, where $R_{\min} := \max\{\varphi'_{\text{ext}} = 0\} \leq R$. In addition, the potential Φ_e is convex.

Proof. — The existence of R is clear. We may have non uniqueness only in the case where Φ_{ext} is constant on a ball $B(0, R_0)$ (of positive radius), since φ'_{ext} is nondecreasing. The potential Φ_e may be searched for under the form of a radial function, and we find the expressions

$$\Phi_e(x) = (\varphi_{\text{ext}}(R) - \varphi_{\text{ext}}(|x|) + \varphi'_{\text{ext}}(R)\Gamma(R))\mathbf{1}_{B(0,R)} + \varphi'_{\text{ext}}(R)\Gamma(x)\mathbf{1}_{\mathbb{R}^N \setminus B(0,R)},$$

where $\Gamma(R)$ stands for $\Gamma(y)$ for any $y \in \partial B(0, R)$. □

Let us give some examples illustrating Proposition 2.16.

EXAMPLE 1 (RADIAL). — If φ_{ext} is of class C^2 and φ''_{ext} is positive on \mathbb{R}_+ , then $n_e(x) = \mathbf{1}_{B(0,R)}(x) \Delta \Phi_{\text{ext}}(x)$ and is absolutely continuous with respect to the Lebesgue measure. We then have $\text{Supp}(n_e) = \overline{B}(0, R)$.

EXAMPLE 2 (RADIAL). — The potential φ_{ext} is C^1 , piecewise C^2 , but is constant on the interval $[0, R_0]$ (hence it is not strictly convex). It does not matter whether the second order derivative of φ_{ext} is continuous or not at R_0 . We define $R \geq R_0 > 0$ by the relation $\mathbf{m} = \int_{B(0,R)} \Delta \Phi_{\text{ext}} \, dx$, or, equivalently, $N|B(0,1)|R^{N-1} \varphi'_{\text{ext}}(R) = \mathbf{m}$. Then, $n_e = \mathbf{1}_{B(0,R) \setminus B(0,R_0)} \Delta \Phi_{\text{ext}}$,

$$\text{Supp}(n_e) = \overline{B}(0, R) \setminus B(0, R_0) \subsetneq \overline{B}(0, R) = \{\Phi_e = 0\}$$

and this set is then neither star-shaped nor simply connected. Here again, if $\varphi_{\text{ext}} \in C^2$, this potential does not fit hypothesis (H2) since $\Delta \Phi_{\text{ext}}$ is not bounded away from 0 near R_0 .

Let us give now examples which do not fit the regularity hypotheses required in Proposition 2.15. These expressions are justified through the characterization (iii) in Theorem 2.7 and simple computation of $\Gamma \star n_e$.

EXAMPLE 3 (RADIAL). — Take the potential $\varphi_{\text{ext}}(r) = r$, that is $\Phi_{\text{ext}}(x) = |x|$. Then, $\Delta\Phi_{\text{ext}} = (N-1)/r > 0$, $n_e = (N-1)|x|^{-1}\mathbf{1}_{B(0,R)}$, with $N|B(0,1)|R^{N-1} = \mathbf{m}$, and $\text{Supp}(n_e) = \overline{B}(0, R) = \{\Phi_e = 0\}$.

EXAMPLE 5 (RADIAL). — Consider the potential $\varphi_{\text{ext}}(r) = r + \max(r-1, 0)$:

– if $\mathbf{m} \leq N|B(0,1)|$, then

$$n_e = (N-1)|x|^{-1}\mathbf{1}_{B(0,R)}, \quad \text{with } R = (\mathbf{m}/N|B(0,1)|)^{1/N-1}$$

and $\text{Supp}(n_e) = \overline{B}(0, R) = \{\Phi_e = 0\}$;

– if $N|B(0,1)| \leq \mathbf{m} \leq 2N|B(0,1)|$, then

$$n_e = (N-1)|x|^{-1}\mathbf{1}_{B(0,1)} + (\mathbf{m} - N|B(0,1)|)\delta_{\partial B(0,1)}$$

and $\text{Supp}(n_e) = \overline{B}(0, 1) = \{\Phi_e = 0\}$;

– if $\mathbf{m} \geq 2N|B(0,1)|$, then

$$n_e = (N-1)|x|^{-1}\mathbf{1}_{B(0,1)} + N|B(0,1)|\delta_{\partial B(0,1)} + 2(N-1)|x|^{-1}\mathbf{1}_{B(0,R) \setminus B(0,1)},$$

where $R \geq 1$ is such that

$$2N|B(0,1)|(R^{N-1} - 1) + N|B(0,1)| = \mathbf{m},$$

and $\text{Supp}(n_e) = \overline{B}(0, R) = \{\Phi_e = 0\}$.

Since we assume φ_{ext} convex and with 0 as a minimum point, it follows that φ_{ext} has a right-derivative at 0, hence the singularity in $1/|x|$ at the origin for n_e is the worst we can have. The radial Example 5 also shows that we may have Dirac masses on a sphere (of positive radius).

Our next results guarantees that \mathcal{K} has non empty interior when the confining potential Φ_{ext} is C^1 and convex.

PROPOSITION 2.17. — *We assume that 0 is a minimum point of Φ_{ext} and that the potential Φ_{ext} is of class C^1 and convex. Then, there exists $r_0 > 0$ such that $B(0, r_0) \subset \mathcal{K}$. In particular, \mathcal{K} has non empty interior.*

Proof. — We follow the argument of [30, Chap. 5, Th. 6.2], where we work on $h := \Gamma \star n_e$ and shall use that it is a solution to the obstacle problem given in Proposition 2.9 with the obstacle $\psi = C_* - \Phi_{\text{ext}}$.

We first consider the case $N \geq 3$ and notice that $\text{Supp}(n_e)$ has a positive capacity: we fix some $a \in \text{Supp}(n_e)$ such that $C_* = h(a) + \Phi_{\text{ext}}(a)$ (see (22)). Now, since $N \geq 3$, we observe that (with $c_N > 0$)

$$h(a) = \Gamma \star n_e(a) = c_N |\cdot|^{2-N} \star n_e > 0$$

and that 0 is actually a global minimum point of Φ_{ext} , thus $C_* > \Phi_{\text{ext}}(a) \geq \Phi_{\text{ext}}(0)$ and it follows that $\psi(0) = C_* - \Phi_{\text{ext}}(0) > 0$. On the other hand, $h(x) \sim \mathbf{m}\Gamma(x)$ tends

to $0 < \psi(0)$ at infinity, thus there exists an $R_0 > 0$ such that $h(x) \leq \psi(0)/2$ when $|x| \geq R_0$. For x_0 that will be close to 0, we let

$$v(x) := \psi(x_0) + (x - x_0) \cdot \nabla\psi(x_0)$$

be the affine tangent to ψ at x_0 . Since ψ is concave (Φ_{ext} is convex), we have $\psi \leq v$ in \mathbb{R}^N . Furthermore, if x_0 is sufficiently close to 0 (depending on R_0), then $\nabla\psi(x_0)$ is small (since ψ is C^1 and achieves a minimum at 0) and thus $v > \psi(0)/2 > 0$ on $\partial B(0, R_0)$. Since $\Delta v \equiv 0$, we may now apply [30, Chap. 4, Th. 8.3] to infer $h \leq v$ in $B(0, R_0)$ (this is a maximum type principle proved using the comparison function

$$g := \min(h, v)\mathbf{1}_{B(0, R_0)} + h\mathbf{1}_{\mathbb{R}^N \setminus B(0, R_0)}$$

in the formulation of the obstacle problem given in Proposition 2.9). In particular, $\psi(x_0) \leq h(x_0) \leq v(x_0) = \psi(x_0)$, which means that, as wished, $x_0 \in \mathcal{K}$.

Let us now turn to the dimensions $N = 2$ and $N = 1$. Then, it may happen that $\psi(0) \leq 0$, but since $h(x) \sim m\Gamma(x)$ tends to $-\infty < \psi(0)$ at infinity, the previous argument still applies. \square

If one is able to prove that \mathcal{K} is convex and assuming that Φ_{ext} satisfies (H2), then (H1) is automatically true. Indeed, any point of $\partial\mathcal{K}$ has then a positive density and we may then apply the regularity result of L. Caffarelli (see e.g. [16, Chap. 2, Th. 3.10]) which ensures that $\partial\Omega$ is of class C^1 (hence C^{s+1} by (H2)).

We conclude with a result from [26, Th. 3.24] on the topology of \mathcal{K} valid only in space dimension two (the proof uses complex analysis).

PROPOSITION 2.18 ([26]). — *We assume $N = 2$. Suppose that Φ_{ext} is of class C^2 and that its Hessian is everywhere positive definite. Then, $\text{Supp}(n_e)$ is simply connected, and equal to the closure of its interior. Moreover, if Φ_{ext} is $C^{2,\alpha}$ for some $\alpha \in]0, 1[$, then $\partial\mathcal{K}$ is a $C^{1,\beta}$ Jordan curve, for some $\beta \in]0, 1[$.*

The above result does not prevent cusps in $\partial\mathcal{K}$, but just says that the boundary $\partial\mathcal{K}$ possesses a $C^{1,\beta}$ parametrization.

3. ASYMPTOTIC ANALYSIS

This section is devoted to the analysis of the asymptotic regime $\varepsilon \rightarrow 0$. We shall point out the difficulties and necessary adaptations between the case of quadratic potentials, Theorem 1.2 and Theorem 2.6, and the general case, Theorem 2.14. For the existence theory of the Vlasov–Poisson equation, we refer the reader to [1] for weak solutions and more recently to [34, 41] where strong solutions and regularity issues are discussed. Further details and references can be found in the survey [20].

3.1. A USEFUL ESTIMATE ON Φ_e . — Before we turn to the analysis of the asymptotic regime $\varepsilon \rightarrow 0$, it is convenient to set up an estimate that describes the behavior of Φ_e close to the neighborhood of $\partial\mathcal{K}$. In the isotropic case, Φ_{ext} being given by (2), the potential Φ_e is defined by (9), and we observe that there exists $C > 0$ such that

$$(28) \quad 0 \leq (|x| - R) |\nabla_x \Phi_e(x)| \leq C \Phi_e(x)$$

holds for any x with $|x| \geq R$. More generally, for a quadratic potential (12), we can establish the following property, based on the formulas in Section 2.1.

LEMMA 3.1. — *Let Φ_e be the quadratic potential defined as in Corollary 2.5. Let $\mathcal{V} : \mathbb{R}^N \rightarrow \mathbb{R}^N$ be smooth, compactly supported and such that $\mathcal{V} \cdot \nu|_{\partial\mathcal{K}_a} = 0$. Then, there exists a positive constant C , depending only on N , Φ_{ext} and \mathcal{V} such that we have, for any $x \in \mathbb{R}^N$,*

$$(29) \quad |\mathcal{V} \cdot \nabla\Phi_e(x)| \leq C\Phi_e(x).$$

Proof. — Since \mathcal{V} is compactly supported and Φ_e is positive in $\{\sigma_a > 0\}$, we just need to prove the inequality for x close to $\partial\mathcal{K}_a$, that is for $\sigma_a(x)$ small. We still define $\lambda > 0$ so that $\lambda^{-2} = \sum_{j=1}^N \lambda_j^{-2}$. From (20), and by Taylor expansion of the integral, we infer that for $0 < \sigma_a(x) \ll 1$ and $1 \leq k \leq N$,

$$\lambda^2 \partial_k \Phi_e(x) = \frac{x_k}{2} \left(\prod_{j=1}^N a_j \right) \left(\sigma_a(x) \frac{1}{a_k^2} \left(\prod_{j=1}^N a_j^2 \right)^{-1/2} + \mathcal{O}(\sigma_a^2(x)) \right) = \frac{x_k \sigma_a(x)}{2a_k^2} + \mathcal{O}(\sigma_a^2(x)).$$

Let $\mathcal{X}(x)$ stands for the vector with components x_k/a_k^2 . In particular, for $0 < \sigma_a(x) \ll 1$, we get

$$(30) \quad |\nabla\Phi_e(x)| = \mathcal{O}(\sigma_a(x)) \quad \text{and} \quad \frac{\nabla\Phi_e(x)}{|\nabla\Phi_e(x)|} = \frac{\mathcal{X}(x)}{|\mathcal{X}(x)|} + \mathcal{O}(\sigma_a(x)),$$

where the unit vector field $x \mapsto \mathcal{X}(x)/|\mathcal{X}(x)|$ is smooth near $\partial\mathcal{K}_a$ and is the (outward) normal on $\partial\mathcal{K}_a$. Now, observe that

$$\begin{aligned} 0 &= \partial_k \left(\sum_{j=1}^N \frac{x_j^2}{a_j^2 + \sigma_a(x)} \right) = 2 \frac{x_k}{a_k^2 + \sigma_a(x)} - \left(\sum_{j=1}^N \frac{x_j^2}{(a_j^2 + \sigma_a(x))^2} \right) \partial_k \sigma_a(x) \\ &= 2 \frac{x_k}{a_k^2} - \partial_k \sigma_a(x) \left(\sum_{j=1}^N \frac{x_j^2}{a_j^4} \right) + \mathcal{O}(\sigma_a(x)). \end{aligned}$$

Therefore, for $0 < \sigma_a(x) \ll 1$ and $1 \leq k \leq N$, we have

$$\begin{aligned} \lambda^2 \partial_k \Phi_e(x) &= \frac{1}{4} \sigma_a(x) \partial_k \sigma_a(x) \left(\sum_{j=1}^N \frac{x_j^2}{a_j^4} \right) + \mathcal{O}(\sigma_a^2(x)) \\ &= \frac{1}{8} \partial_k \left(\sigma_a^2(x) \sum_{j=1}^N \frac{x_j^2}{a_j^4} \right) + \mathcal{O}(\sigma_a^2(x)). \end{aligned}$$

As a consequence,

$$(31) \quad \lambda^2 \Phi_e(x) = \frac{1}{8} \sigma_a^2(x) \left(\sum_{j=1}^N \frac{x_j^2}{a_j^4} \right) + \mathcal{O}(\sigma_a^3(x)) \geq \frac{\sigma_a^2(x)}{C},$$

holds for some $C > 0$. Going back to (30), we arrive at

$$\begin{aligned} \mathcal{V}(x) \cdot \nabla \Phi_e(x) &= \mathcal{V}(x) \cdot \left(\frac{\nabla \Phi_e(x)}{|\nabla \Phi_e(x)|} \right) \times |\nabla \Phi_e(x)| \\ &= \mathcal{V}(x) \cdot \left(\frac{\mathcal{X}(x)}{|\mathcal{X}(x)|} + \mathcal{O}(\sigma_a(x)) \right) \times \mathcal{O}(\sigma_a(x)) \\ &= (\mathcal{O}(\sigma_a(x)) + \mathcal{O}(\sigma_a(x))) \times \mathcal{O}(\sigma_a(x)) = \mathcal{O}(\sigma_a^2(x)) = \mathcal{O}(\Phi_e(x)), \end{aligned}$$

by (31) and since $\mathcal{V} \cdot \mathcal{X}/|\mathcal{X}|$ vanishes when $\sigma_a = 0$ in view of the no-flux condition satisfied by \mathcal{V} . This finishes the proof. \square

In the more general setting considered in Theorem 2.14, the result is the following and simply relies on the use of a local chart.

LEMMA 3.2. — We assume that $\partial\Omega$ is of class C^1 and that (H2) is satisfied. Then, there exists a constant C such that, for any $x \in \mathbb{R}^N$,

$$(32) \quad |\mathcal{V} \cdot \nabla \Phi_e(x)| \leq C \Phi_e(x).$$

Proof. — We have already seen that $\partial\Omega$ is actually of class C^{s+1} . Since Φ_e is positive in $\mathbb{R}^N \setminus \mathcal{K}$ and \mathcal{V} has compact support, by a compactness argument, it suffices to show that (32) holds near any point $a \in \partial\Omega$. Possibly translating and rotating the axis, we assume $a = 0$ and that the inward normal to Ω at $a = 0$ is $e_1 = (1, 0, \dots, 0)$. We let $x_1 = \Theta(x_\perp)$, where $x_\perp = (x_2, \dots, x_N)$, be a C^2 parametrization of $\partial\Omega$ near 0, with $\nabla\Theta(0) = 0$, hence $\Theta(x_\perp) = \mathcal{O}(|x_\perp|^2)$.

We now consider the function $\varphi : \mathbb{R}^N \rightarrow \mathbb{R}$ defined by $\varphi(y) := \Phi_e(y_1 + \Theta(y_\perp), y_\perp)$, where $y_\perp = (y_2, \dots, y_N) \in \mathbb{R}^{N-1}$. Then, $\varphi(y) = 0$ when $y_1 \geq 0$, hence, for $2 \leq j \leq N$ and $1 \leq k \leq N$ and if $y_1 = 0$, $\partial_k \varphi(y) = \partial_{j,k}^2 \varphi(y) = 0$; moreover, $\partial_{1,1}^2 \varphi(0, y_\perp) = \Delta \Phi_{\text{ext}}(0, y_\perp)$ in view of the equality

$$\Delta \Phi_{\text{ext}}(x) = \Delta_x \Phi_e(x) = (\Delta_y \varphi - (\Delta_\perp \Theta) \partial_1 \varphi + \sum_{j=2}^N (\partial_j \Theta)^2 \partial_{1,j}^2 \varphi)(x_1 - \Theta(x_\perp), x_\perp)$$

in $\{x_1 \leq \Theta(x_\perp)\}$.

It follows from these relations that, by the Taylor formula and by using $\Delta \Phi_{\text{ext}}(0) > 0$ and $y_1 = x_1 - \Theta(x_\perp) \leq 0$,

$$\varphi(y) = \varphi(y) - \varphi(0, y_\perp) - y_1 \partial_1 \varphi(0, y_\perp) = y_1^2 \int_0^1 (1-t) \partial_1^2 \varphi(ty_1, y_\perp) dt \geq \frac{y_1^2}{C},$$

and we deduce

$$(33) \quad \Phi_e(x) \geq \frac{(x_1 - \Theta(x_\perp))^2}{C}.$$

Still by the Taylor formula, we have, for $2 \leq j \leq N$,

$$\partial_j \varphi(y) = y_1^2 \int_0^1 (1-t) \partial_{1,1,j}^3 \varphi(ty_1, y_\perp) dt = \mathcal{O}(y_1^2)$$

and

$$\partial_1 \varphi(y) = y_1 \partial_1^2 \varphi(0, y_\perp) + y_1^2 \int_0^1 (1-t) \partial_1^3 \varphi(ty_1, y_\perp) dt = y_1 \Delta \Phi_{\text{ext}}(0, y_\perp) + \mathcal{O}(y_1^2).$$

Now, we write $\partial_1 \Phi_e(x) = \partial_1 \varphi(y)$ (with $y = (x_1 - \Theta(x_\perp), x_\perp)$) and $\nabla_\perp \Phi_e(x) = \nabla_\perp \varphi(y) - \partial_1 \varphi(y) \nabla_\perp \Theta(y_\perp)$, thus

$$\begin{aligned} \mathcal{V}(x) \cdot \nabla \Phi_e(x) &= \mathcal{V}_1(x) \partial_1 \Phi_e(x) + \mathcal{V}_\perp(x) \cdot \nabla_\perp \Phi_e(x) \\ &= \mathcal{V}_1(x) \partial_1 \varphi(y) + \mathcal{V}_\perp(x) \cdot \nabla_\perp \varphi(y) - \partial_1 \varphi(y) \mathcal{V}_\perp(x) \cdot \nabla_\perp \Theta(y_\perp). \end{aligned}$$

Note that $\nabla_\perp \varphi(y) = \mathcal{O}(y_1^2)$. Furthermore, since $\mathcal{V} \cdot \nu = 0$ on $\partial\Omega = \{x_1 = \Theta(x_\perp)\}$ and $\nu(x) = (1, -\nabla_\perp \Theta(x_\perp)) / |(1, -\nabla_\perp \Theta(x_\perp))|$, we deduce

$$\begin{aligned} \mathcal{V}(x) \cdot \nabla \Phi_e(x) &= \mathcal{O}(y_1^2) + \partial_1 \varphi(y) \left([\mathcal{V}_1(x) - \mathcal{V}_\perp(x) \cdot \nabla_\perp \Theta(x_\perp)] \right. \\ &\quad \left. - [\mathcal{V}_1(\Theta(x_\perp), x_\perp) - \mathcal{V}_\perp(\Theta(x_\perp), x_\perp) \cdot \nabla_\perp \Theta(x_\perp)] \right) \\ &= \mathcal{O}(y_1^2) + \mathcal{O}(|y_1|) \times \mathcal{O}(|x_1 - \Theta(x_\perp)|) = \mathcal{O}(y_1^2). \end{aligned}$$

We conclude by using (33). \square

3.2. BASIC A PRIORI ESTIMATES. — Now that we have in hand the limiting density profile n_e and the associated potential field Φ_e , we derive some basic a priori estimates from (V)–(P).

Using the splitting of Poisson equation as in (10), (V) recasts as

$$\partial_t f_\varepsilon + v \cdot \nabla_x f_\varepsilon - \frac{1}{\varepsilon} \nabla_x \Phi_e \cdot \nabla_v f_\varepsilon - \frac{1}{\sqrt{\varepsilon}} \nabla_x \Psi_\varepsilon \cdot \nabla_v f_\varepsilon = 0.$$

Let us compute the time variation of the following energies:

– Kinetic energy

$$\frac{d}{dt} \iint \frac{|v|^2}{2} f_\varepsilon dv dx = -\frac{1}{\varepsilon} \iint v \cdot \nabla_x \Phi_e f_\varepsilon dv dx - \frac{1}{\sqrt{\varepsilon}} \iint v \cdot \nabla_x \Psi_\varepsilon f_\varepsilon dv dx.$$

– Leading order potential energy

$$\frac{d}{dt} \iint \Phi_e f_\varepsilon dv dx = \iint v \cdot \nabla_x \Phi_e f_\varepsilon dv dx.$$

– Fluctuations potential energy

$$\begin{aligned} \frac{d}{dt} \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx &= \int \nabla_x \Psi_\varepsilon \cdot \partial_t \nabla_x \Psi_\varepsilon dx = - \int \Psi_\varepsilon \partial_t \left(\frac{n_e - \rho_\varepsilon}{\sqrt{\varepsilon}} \right) dx \\ &= - \int \Psi_\varepsilon \frac{1}{\sqrt{\varepsilon}} \nabla_x \cdot \left(\int v f_\varepsilon dv \right) dx = \frac{1}{\sqrt{\varepsilon}} \iint v \cdot \nabla_x \Psi_\varepsilon f_\varepsilon dv dx. \end{aligned}$$

By summing these relations, we conclude with the following claim (which applies for all three cases for Φ_{ext}).

PROPOSITION 3.3. — *The solution $(f_\varepsilon, \Phi_\varepsilon = \frac{1}{\varepsilon} \Phi_e + \frac{1}{\sqrt{\varepsilon}} \Psi_\varepsilon)$ of (V)–(P) satisfies the following energy conservation equality*

$$\frac{d}{dt} \left\{ \iint \frac{|v|^2}{2} f_\varepsilon dv dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon dv dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx \right\} = 0.$$

Furthermore, the total charge is conserved

$$\iint f_\varepsilon(t, x, v) \, dv \, dx = \iint f_\varepsilon(0, x, v) \, dv \, dx = \mathbf{m}.$$

3.3. CONVERGENCE OF THE DENSITY AND THE CURRENT. — We assume a uniform bound on the energy at the initial time, namely

$$(34) \quad \sup_{0 < \varepsilon < 1} \iint \frac{1}{2} |v|^2 f_\varepsilon^{\text{init}} \, dv \, dx + \frac{1}{\varepsilon} \iint \Phi_\varepsilon f_\varepsilon^{\text{init}} \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon^{\text{init}}|^2 \, dx < \infty,$$

where $\Psi_\varepsilon^{\text{init}}$ solves the Poisson equation (10). Then, Proposition 3.3 ensures that the energy remains uniformly bounded for positive times. Thus, possibly at the price of extracting subsequences, we can suppose that

$$\begin{aligned} f_\varepsilon &\rightharpoonup f \quad \text{weakly-}\star \text{ in } \mathcal{M}^1([0, T] \times \mathbb{R}^N \times \mathbb{R}^N), \\ \rho_\varepsilon &= \int f_\varepsilon \, dv \rightharpoonup \rho \quad \text{weakly-}\star \text{ in } \mathcal{M}^1([0, T] \times \mathbb{R}^N). \end{aligned}$$

Going back to the Poisson equation, we observe that

$$n_e - \rho_\varepsilon = \sqrt{\varepsilon} \nabla_x \cdot (\nabla_x \Psi_\varepsilon)$$

where, by Proposition 3.3, $\nabla_x \Psi_\varepsilon$ is bounded in $L^\infty(0, T; L^2(\mathbb{R}^N))$. Consequently, we establish the following claim.

LEMMA 3.4. — *The sequence ρ_ε converges to $n_e = \rho$ strongly in $L^\infty(0, T; H^{-1}(\mathbb{R}^N))$ and weakly- \star in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$. The limit f is supported in $[0, T] \times \bar{\Omega} \times \mathbb{R}^N$. The sequence $J^\varepsilon = \int v f_\varepsilon \, dv$ is bounded in $L^\infty(0, T; L^1(\mathbb{R}^N))$; it admits a subsequence which converges, say weakly- \star in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$; the limit J is divergence free, supported in $[0, T] \times \bar{\Omega}$ and may be written $\int v f \, dv = n_e W$ for some $W \in \mathcal{M}^1([0, T] \times \mathbb{R}^N)$.*

Proof. — Proposition 3.3 tells us that $|v|^2 f_\varepsilon$ is bounded in $L^\infty(0, T; L^1(\mathbb{R}^N \times \mathbb{R}^N))$. Hence, by using Cauchy-Schwarz' inequality, we get

$$(35) \quad \begin{aligned} \int |J_\varepsilon| \, dx &\leq \iint |v| \sqrt{f_\varepsilon} \sqrt{f_\varepsilon} \, dv \, dx \\ &\leq \left(\iint |v|^2 f_\varepsilon \, dv \, dx \right)^{1/2} \left(\iint f_\varepsilon \, dv \, dx \right)^{1/2}, \end{aligned}$$

which leads to the asserted uniform estimate on the current. We can thus also assume $J_\varepsilon \rightharpoonup J$ weakly- \star in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$. Furthermore, since the second order moment in v of f_ε is uniformly bounded, we check that

$$\rho = \int f \, dv, \quad J = \int v f \, dv.$$

Note that ρ_ε and J_ε satisfy (5). Letting ε go to 0 yields

$$\partial_t \rho + \nabla_x \cdot J = 0 = \partial_t n_e + \nabla_x \cdot J = 0 + \nabla_x \cdot J = 0.$$

Thus, J is divergence-free. Finally, on the one hand, we have $\rho = \int f \, dv$, and on the other hand $\rho = n_e$ is absolutely continuous with respect to the Lebesgue measure and

supported in $[0, T] \times \bar{\Omega}$. It follows that $\text{Supp}(f) \subset [0, T] \times \bar{\Omega} \times \mathbb{R}^N$, and thus $\text{Supp}(J) \subset [0, T] \times \bar{\Omega}$; furthermore we have $f([0, T] \times \partial\Omega \times \mathbb{R}^N) = 0$, and $J([0, T] \times \partial\Omega) = 0$. \square

In order to define the normal trace of J over $\partial\Omega$ (that is the sphere $\partial B(0, R)$ in the case (2)), we shall use the theory introduced in [9]. As a consequence of the discussion above, we start by observing that J belongs to the set $\mathcal{DM}^{\text{ext}}(\mathbb{R}^N)$ of extended divergence-measure fields over \mathbb{R}^N , see [9, Def. 1.1]. Therefore, according to [9, Th. 3.1], J admits a normal trace $J \cdot \nu|_{\partial\Omega}$ defined as a continuous linear functional over $\text{Lip}(\gamma, \partial\Omega)$, $\gamma > 1$ (see [9, Eq. (2.1)]) with

$$\langle J \cdot \nu|_{\partial\Omega}, \phi \rangle = \int_{\Omega} \hat{\phi} \nabla_x \cdot J + \int_{\Omega} J \cdot \nabla_x \hat{\phi},$$

where the function $\hat{\phi} \in \text{Lip}(\gamma, \Omega)$ in the right-hand side is an extension of $\phi \in \text{Lip}(\gamma, \partial\Omega)$. However, by $\nabla_x \cdot J = 0$ and the support property on J , we can rewrite

$$\langle J \cdot \nu|_{\partial\Omega}, \phi \rangle = 0 + \int_{\mathbb{R}^N} J \cdot \nabla_x \hat{\phi} = -\langle \nabla_x \cdot J, \hat{\phi} \rangle = 0.$$

Another way to see this is to observe that the normal trace from Ω must be the same as the normal trace from $\mathbb{R}^N \setminus \bar{\Omega}$, which is clearly zero since J has support in $\bar{\Omega}$. Consequently,

$$J \cdot \nu|_{\bar{\Omega}} = 0 \quad \text{in } [C(0, T; \text{Lip}(\gamma, \bar{\Omega}))]^*.$$

Remark that this is not a pointwise relation. In particular, it may happen that $J_\varepsilon^{\text{init}} \cdot \nu|_{\bar{\Omega}} = \rho_\varepsilon V_\varepsilon^{\text{init}} \cdot \nu|_{\bar{\Omega}}$ is nonzero, but this does not prevent the time integral of $J \cdot \nu|_{\bar{\Omega}}$ to vanish.

3.4. PASSING TO THE LIMIT: MODULATED ENERGY. — As in [5], we now study the modulated energy

$$\mathcal{H}_{\mathcal{V}, \varepsilon} = \frac{1}{2} \iint |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon \, dv \, dx,$$

where all the terms integrated are nonnegative. Let us compute as follows

$$\frac{d}{dt} \mathcal{H}_{\mathcal{V}, \varepsilon} = \frac{d}{dt} \iint \left(\frac{|\mathcal{V}|^2}{2} - \mathcal{V} \cdot v \right) f_\varepsilon \, dv \, dx = \frac{d}{dt} \int \left(\rho_\varepsilon \frac{|\mathcal{V}|^2}{2} - \mathcal{V} \cdot J_\varepsilon \right) dx,$$

by using Proposition 3.3. We thus have

$$\frac{d}{dt} \mathcal{H}_{\mathcal{V}, \varepsilon} = \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot \partial_t \mathcal{V} \, dx + \int \frac{|\mathcal{V}|^2}{2} \partial_t \rho_\varepsilon \, dx - \int \mathcal{V} \cdot \partial_t J_\varepsilon \, dx.$$

Here, we are assuming that the solution f_ε of the Vlasov–Poisson system (V)–(P) is regular enough so that we can perform all the calculations that follow. Integrating the Vlasov equation, we obtain

$$(36) \quad \partial_t J_\varepsilon + \nabla_x \cdot \mathbb{P}_\varepsilon + \frac{1}{\sqrt{\varepsilon}} \rho_\varepsilon \nabla_x \Psi_\varepsilon + \frac{1}{\varepsilon} \rho_\varepsilon \nabla_x \Phi_e = 0,$$

where we rewrite

$$\frac{1}{\sqrt{\varepsilon}} \rho_\varepsilon \nabla_x \Psi_\varepsilon = \frac{\rho_\varepsilon - n_e}{\sqrt{\varepsilon}} \nabla_x \Psi_\varepsilon + \frac{n_e}{\sqrt{\varepsilon}} \nabla_x \Psi_\varepsilon = -\Delta_x \Psi_\varepsilon \nabla_x \Psi_\varepsilon + \frac{n_e}{\sqrt{\varepsilon}} \nabla_x \Psi_\varepsilon,$$

and

$$\Delta_x \Psi_\varepsilon \nabla_x \Psi_\varepsilon = \nabla_x \cdot (\nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) - \nabla_x (|\nabla_x \Psi_\varepsilon|^2/2).$$

Combining these relations to the charge conservation (5) and integration by parts, we arrive at

$$\begin{aligned} \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon} &= \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot \partial_t \mathcal{V} \, dx + \int J_\varepsilon \cdot \nabla_x (|\mathcal{V}|^2/2) \, dx \\ &\quad - \int D_x \mathcal{V} : (\mathbb{P}_\varepsilon - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) \, dx + \int \nabla_x \cdot \mathcal{V} \frac{|\nabla_x \Psi_\varepsilon|^2}{2} \, dx \\ &\quad + \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_e \, dx + \int \frac{n_e}{\sqrt{\varepsilon}} \mathcal{V} \cdot \nabla_x \Psi_\varepsilon \, dx, \end{aligned}$$

where $D_x \mathcal{V}$ stands for the jacobian matrix of the vector field \mathcal{V} . For the last integral, since n_e is supported in Ω , we write it as

$$\int_\Omega \frac{n_e}{\sqrt{\varepsilon}} \mathcal{V} \cdot \nabla_x \Psi_\varepsilon \, dx = 0$$

by integration by parts and using that $\nabla_x \cdot (n_e \mathcal{V}) = 0$ in Ω and the no-flux condition (6).

Let us set

$$\mathbb{P}_{\mathcal{V},\varepsilon} := \int (v - \mathcal{V}) \otimes (v - \mathcal{V}) f_\varepsilon \, dv = \mathbb{P}_\varepsilon - \mathcal{V} \otimes J_\varepsilon - J_\varepsilon \otimes \mathcal{V} + \rho_\varepsilon \mathcal{V} \otimes \mathcal{V}.$$

A direct substitution leads to

$$\begin{aligned} (37) \quad \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon} &= \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V}) \, dx \\ &\quad - \int D_x \mathcal{V} : (\mathbb{P}_{\mathcal{V},\varepsilon} - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) \, dx \\ &\quad + \int \nabla_x \cdot \mathcal{V} \frac{|\nabla_x \Psi_\varepsilon|^2}{2} \, dx + \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_e \, dx. \end{aligned}$$

We shall use the shorthand notation $A \lesssim B$ when the inequality $A \leq CB$ holds for some constant $C > 0$, the value of which might vary from a line to another. As a matter of fact, we can dominate the second and third integrals of the right-hand side by

$$\|D_x \mathcal{V}\|_\infty \left(\iint |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx + \int |\nabla_x \Psi_\varepsilon|^2 \, dx \right) \leq \|D_x \mathcal{V}\|_\infty \mathcal{H}_{\mathcal{V},\varepsilon}.$$

Let us distinguish the case of the isotropic potential in order to point out the difficulties. When Φ_{ext} is given by (2), we remind the reader that Φ_e is supported in $\{|x| \geq R\}$, radially symmetric and increasing in $|x|$, see (9). Combining this with (28) allows us to estimate the last term in (37) as follows:

$$\begin{aligned} \left| \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_e \, dx \right| &= \left| \frac{1}{\varepsilon} \int_{\{|x|>R\}} \rho_\varepsilon \frac{\mathcal{V} \cdot x/|x|}{|x| - R} (|x| - R) |\nabla_x \Phi_e| \, dx \right| \\ &\lesssim \frac{1}{\varepsilon} \int_{\{|x|>R\}} \rho_\varepsilon \Phi_e \, dx \left\| \frac{\mathcal{V} \cdot x/|x|}{|x| - R} \right\|_\infty \lesssim \mathcal{H}_{\mathcal{V},\varepsilon}, \end{aligned}$$

where we have used that $\frac{\mathcal{V} \cdot x/|x|}{|x|-R}$ belongs to $L^\infty(\mathbb{R}^N)$ since \mathcal{V} is smooth, compactly supported, and $\mathcal{V} \cdot \nu = 0$ on $\partial B(0, R)$. For a quadratic external potential (12), we can proceed similarly by using Lemma 3.1. When dealing with a general potential, we made hypothesis (H2) so that Lemma 3.2 applies and (29) allows us to estimate

$$\left| \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_e dx \right| \lesssim \frac{1}{\varepsilon} \int \rho_\varepsilon \Phi_e dx \lesssim \mathcal{H}_{\mathcal{V}, \varepsilon}.$$

Therefore, we obtain

$$(38) \quad \frac{d}{dt} \mathcal{H}_{\mathcal{V}, \varepsilon} \lesssim \mathcal{H}_{\mathcal{V}, \varepsilon} + r_\varepsilon$$

where we have set

$$r_\varepsilon := \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V}) dx.$$

The Grönwall lemma yields

$$\mathcal{H}_{\mathcal{V}, \varepsilon}(t) \leq e^{Ct} \left(\mathcal{H}_{\mathcal{V}, \varepsilon}(0) + \int_0^t e^{-C\tau} r_\varepsilon(\tau) d\tau \right),$$

for a certain constant $C > 0$. The assumption (11) on the initial data is that $\lim_{\varepsilon \rightarrow 0} \mathcal{H}_{\mathcal{V}, \varepsilon}(0) = 0$. Hence, we are left with the task of proving that $\int_0^t r_\varepsilon(\tau) d\tau$ tends to 0 as $\varepsilon \rightarrow 0$. We have

$$\begin{aligned} \int_0^t r_\varepsilon(\tau) d\tau &\xrightarrow{\varepsilon \rightarrow 0} \int_0^t \int (n_e \mathcal{V} - J) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V}) dx d\tau \\ &= \int_0^t \int_\Omega (n_e \mathcal{V} - J) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V}) dx d\tau \\ &= - \int_0^t \int_\Omega (n_e \mathcal{V} - J) \cdot \nabla_x p dx d\tau = 0, \end{aligned}$$

since $n_e \mathcal{V}$ and J are divergence free on Ω and their normal trace vanish. \square

It is worth pointing out that the regularity assumption of the sequence of solutions f_ε was only made to justify the computations leading to (38). If one consider less regular solutions, we have to assume that these solutions were constructed through a regularization procedure and that the previous calculations were done on these regularizations and hence (38) will still hold.

3.5. IDENTIFICATION OF THE LIMIT. — Let us observe that if the initial datum satisfies (11), then (34) holds true. Let us first justify (i): we shall show that $\int \rho_\varepsilon \chi dx \rightarrow \int n_e \chi dx$ uniformly on $[0, T]$ as $\varepsilon \rightarrow 0$ for any $\chi \in C_0^0(\mathbb{R}^N)$. We start by observing that

$$(39) \quad \left| \int \rho_\varepsilon(t, x) \chi(x) dx \right| \leq \mathfrak{m} \|\chi\|_\infty$$

holds for any $\chi \in C_0^0(\mathbb{R}^N)$. Next, consider $\chi \in C_c^1(\mathbb{R}^N)$. The charge conservation (5) yields

$$\frac{d}{dt} \int \rho_\varepsilon(t, x) \chi(x) dx = \int \partial_t \rho_\varepsilon \chi dx = - \int \nabla_x \cdot J_\varepsilon \chi dx = \int J_\varepsilon \cdot \nabla_x \chi dx,$$

hence the uniform bound (35) on J_ε implies a uniform bound on $\frac{d}{dt} \int \rho_\varepsilon(t, x)\chi(x) dx$ for $0 \leq t \leq T$. By virtue of the Ascoli-Arzelà theorem, the set

$$\left\{ t \mapsto \int \rho_\varepsilon(t, x)\chi(x) dx, \varepsilon > 0 \right\}$$

is therefore relatively compact in $C([0, T])$ for any fixed $\chi \in C_c^1(\mathbb{R}^N)$. This property extends to any $\chi \in C_0^0(\mathbb{R}^N)$ by virtue of (39). Indeed, for any $\delta > 0$, we can pick $\chi_\delta \in C_c^1(\mathbb{R}^N)$ such that $\|\chi - \chi_\delta\|_\infty \leq \delta/m$. It follows that

$$\int \rho_\varepsilon(t, x)\chi(x) dx = \int \rho_\varepsilon(t, x)(\chi - \chi_\delta)(x) dx + \int \rho_\varepsilon(t, x)\chi_\delta(x) dx,$$

where, owing to (39), the former integral is uniformly dominated by δ and the latter lies in a compact set of $C([0, T])$. Therefore $\{t \mapsto \int \rho_\varepsilon(t, x)\chi(x) dx, \varepsilon > 0\}$ can be covered by a finite number of balls with radius 2δ in $C([0, T])$. Finally, since $C_0^0(\mathbb{R}^N)$ is separable, we apply a diagonal argument to extract a subsequence such that $\int \rho_\varepsilon(t, x)\chi(x) dx$ converges uniformly in $C([0, T])$ for any element χ of a numerable dense set in $C_0^0(\mathbb{R}^N)$. By uniqueness of the limit, we find

$$\lim_{\varepsilon \rightarrow 0} \int \rho_\varepsilon(t, x)\chi(x) dx = \int n_e \chi dx.$$

Going back to (39), we check that the convergence holds for any $\chi \in C_0^0(\mathbb{R}^N)$.

The manipulations detailed in the previous section prove (ii). According to [4, 5], in order to establish (iii), we introduce the following functional: given λ a non negative bounded measure on $[0, T] \times \mathbb{R}^N$, and μ a vector valued bounded measure on $[0, T] \times \mathbb{R}^N$, we set

$$\mathcal{H}(\lambda, \mu) := \sup_{\Theta} \left\{ \int \mu \cdot \Theta - \frac{1}{2} \int \lambda |\Theta|^2 \right\},$$

where the supremum is taken over continuous functions $\Theta : [0, T] \times \mathbb{R}^N \rightarrow \mathbb{R}^N$. According to [4, Prop. 3.4], we have:

LEMMA 3.5 ([4]). — *If μ is absolutely continuous with respect to λ , denoting by \mathbb{V} the Radon-Nikodym derivative of μ with respect to λ , we have*

$$\mathcal{H}(\lambda, \mu) = \frac{1}{2} \int \lambda |\mathbb{V}|^2 \in [0, \infty],$$

otherwise $\mathcal{H}(\lambda, \mu) = +\infty$.

Clearly $(\lambda, \mu) \mapsto \mathcal{H}(\lambda, \mu)$ is a convex and lower semi-continuous (for the weak- \star convergence) functional. Let $\eta : [0, T] \rightarrow [0, \infty)$ be a continuous non negative function.

Reasoning as in [5], we show that $J \in L^\infty(0, T; L^2(\mathbb{R}^N))$ since

$$\begin{aligned} \mathcal{K}(\eta\rho_\varepsilon, \eta J_\varepsilon) &= \frac{1}{2} \int_0^T \int_{\mathbb{R}^N} \frac{|J_\varepsilon(t, x)|^2}{\rho_\varepsilon(t, x)} \eta(t) \, dx \, dt \\ &= \frac{1}{2} \int_0^T \int_{\mathbb{R}^N} \frac{1}{\rho_\varepsilon(t, x)} \left| \int_{\mathbb{R}^N} v \sqrt{f_\varepsilon(t, x, v)} \sqrt{f_\varepsilon(t, x, v)} \, dv \right|^2 \eta(t) \, dx \, dt \\ &\leq \frac{1}{2} \int_0^T \iint_{\mathbb{R}^N \times \mathbb{R}^N} |v|^2 f_\varepsilon(t, x, v) \eta(t) \, dv \, dx \, dt \lesssim \int_0^T \eta \, dt \end{aligned}$$

becomes, as ε tends to 0

$$\mathcal{K}(\eta n_e, \eta J) \lesssim \|\eta\|_{L^1(0, T)}.$$

Reasoning the same way, we get

$$\begin{aligned} \mathcal{K}(\rho_\varepsilon, J_\varepsilon - \rho_\varepsilon \mathcal{V}) &= \frac{1}{2} \int_0^T \int_{\mathbb{R}^N} \frac{|J_\varepsilon - \rho_\varepsilon \mathcal{V}|^2}{\rho_\varepsilon} \, dx \, dt \\ &\leq \frac{1}{2} \int_0^T \iint_{\mathbb{R}^N \times \mathbb{R}^N} |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx \, dt \leq \int_0^T \mathcal{H}_{\mathcal{V}, \varepsilon} \, dt. \end{aligned}$$

It follows that $\mathcal{K}(n_e, J - n_e \mathcal{V}) = 0$, which identifies the limit J and ends the proof of (iii).

Finally, we can check that the initial data for the limit equation is meaningful by establishing some time-compactness on the sequence J_ε . Let

$$\mathscr{W}_R = \{\Theta : [0, T] \times \mathbb{R}^N \rightarrow \mathbb{R}^N, \Theta \text{ of class } C^1, \text{Supp}(\Theta) \subset [0, T] \times \bar{\Omega}, \nabla_x \cdot (n_e \Theta) = 0\},$$

which is a closed subspace of the Banach space C^1 (endowed with the sup norm for the function and its first order derivatives). Multiplying (36) by a function in \mathscr{W}_R , we shall get rid of the stiff terms. Indeed, for such a trial function Θ , we deduce from (36)

$$(40) \quad \frac{d}{dt} \int J_\varepsilon \cdot \Theta \, dx = \int J_\varepsilon \cdot \partial_t \Theta \, dx - \int \Theta \cdot (\nabla_x \cdot \mathbb{P}_\varepsilon) \, dx - \frac{1}{\sqrt{\varepsilon}} \int \rho_\varepsilon \Theta \cdot \nabla_x \Psi_\varepsilon \, dx,$$

since $\Theta \cdot \nabla_x \Phi_e = 0$ pointwise in view of the supports. By using the estimates deduced from Proposition 3.3, we observe that the first two terms are bounded in $L^\infty(0, T)$.

For the last one, we use the Poisson equation (10) and integration by parts to infer

$$\begin{aligned} \frac{1}{\sqrt{\varepsilon}} \int \rho_\varepsilon \Theta \cdot \nabla_x \Psi_\varepsilon \, dx &= \frac{1}{\sqrt{\varepsilon}} \int n_e \Theta \cdot \nabla_x \Psi_\varepsilon \, dx - \int \Delta_x \Psi_\varepsilon \Theta \cdot \nabla_x \Psi_\varepsilon \, dx \\ &= 0 + \int \nabla_x \Psi_\varepsilon \cdot \nabla_x (\Theta \cdot \nabla_x \Psi_\varepsilon) \, dx \\ &= \frac{1}{2} \int \Theta \cdot \nabla_x (|\nabla_x \Psi_\varepsilon|^2) \, dx + \sum_{1 \leq j, k \leq N} \int \partial_{x_j} \Psi_\varepsilon \partial_{x_j} \Theta_k \partial_{x_k} \Psi_\varepsilon \, dx, \end{aligned}$$

where we have used that $n_e \Theta(t, \cdot)$ is divergence free. For quadratic external potentials, an integration by parts shows that the first integral is zero (since $n_e \Theta$ is divergence free). In any cases, the right hand side can be dominated

by $\|\nabla\Theta\|_\infty\|\nabla\Psi_\varepsilon\|_{L^\infty(0,T;L^2(\mathbb{R}^N))}$ and it is thus bounded in $L^\infty(0,T)$. Reporting this into (40) allows us to conclude that

$$\frac{d}{dt} \int J_\varepsilon \cdot \Theta \, dx \text{ is bounded in } L^\infty(0,T).$$

Since \mathscr{W}_R is separable, we can boil down a diagonal argument to justify that J_ε is relatively compact in $C^0(0,T;\mathscr{W}'_R\text{-weak-}\star)$: we can assume that the extracted subsequence is such that $\int J_\varepsilon \cdot \Theta \, dx$ converges uniformly on $[0,T]$ for any $\Theta \in \mathscr{W}_R$. \square

4. ASYMPTOTIC ANALYSIS OF THE VLASOV–POISSON–FOKKER–PLANCK SYSTEM

In this section we state and prove a theorem analogous to Theorem 1.2 when the basic equation is (VFP), which includes a Fokker–Planck operator, coupled with (P).

For the well-posedness issues of the system (VFP) coupled to (P), we refer the reader to [3, 12]. The role of the external potential is precisely investigated in [13]. The associated moment system reads

$$\begin{cases} \partial_t \rho_\varepsilon + \nabla_x \cdot J_\varepsilon = 0, \\ \partial_t J_\varepsilon + \nabla_x \cdot \mathbb{P}_\varepsilon + \rho_\varepsilon \nabla_x \Phi_\varepsilon = -J_\varepsilon, \end{cases}$$

where we still use the notation $J_\varepsilon = \int v f_\varepsilon \, dv$, $\mathbb{P}_\varepsilon = \int v \otimes v f_\varepsilon \, dv$. As $\varepsilon \rightarrow 0$, we expect as before that $\rho_\varepsilon \rightarrow n_e = \mathbf{1}_\Omega \Delta \Phi_{\text{ext}}$ and that the behavior of the current is driven by the Lake Equation with friction

$$(LE_f) \quad \begin{cases} \partial_t V + V \cdot \nabla_x V + \nabla_x p = -V, \\ \nabla_x \cdot (n_e V) = 0. \end{cases}$$

If Φ_{ext} is quadratic as in (12) (possibly isotropic), the domain Ω is an ellipsoid (possibly a ball) as in Section 2.1 and (LE_f) becomes the Incompressible Euler system with friction

$$\begin{cases} \partial_t V + \nabla_x \cdot (V \otimes V) + \nabla_x p = -V, \\ \nabla_x \cdot V = 0. \end{cases}$$

For a more general confining potential Φ_{ext} , we make assumptions (h1), (h2), (H1) and (H2) as in Section 2.2. Since we work with finite charge data, the limit equation (LE_f) holds in Ω , completed with the no-flux boundary condition (6), namely

$$V(t, x) \cdot \nu(x) \Big|_{\partial\Omega} = 0.$$

Like in the previous section we associate with V , smooth solution of (LE_f) , a smooth compactly supported extension \mathscr{V} defined on $[0, T] \times \mathbb{R}^N$ such that $\mathscr{V} \cdot \nu(x) \Big|_{\partial\Omega} = 0$.

We shall investigate this asymptotics in the specific case where the “temperature” $\theta = \theta_\varepsilon$ goes to 0 as $\varepsilon \rightarrow 0$. In this context, we can derive an analog of Proposition 3.3 that accounts for the dissipation mechanisms induced by the Fokker–Planck operator.

PROPOSITION 4.1. — *The solution $(f_\varepsilon, \Phi_\varepsilon = \frac{1}{\varepsilon} \Phi_e + \frac{1}{\sqrt{\varepsilon}} \Psi_\varepsilon)$ of (VFP)–(P) satisfies the following entropy dissipation inequality*

$$\frac{d}{dt} \left\{ \frac{1}{2} \iint |v|^2 f_\varepsilon \, dv \, dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon \, dv \, dx + \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx \right\} = -\mathcal{D}_\varepsilon,$$

where we denote

$$\mathcal{D}_\varepsilon = \iint |v\sqrt{f_\varepsilon} + 2\theta_\varepsilon \nabla_v \sqrt{f_\varepsilon}|^2 \, dv \, dx \geq 0.$$

Furthermore, the total charge is conserved

$$\iint f_\varepsilon(t, x, v) \, dv \, dx = \iint f_\varepsilon(0, x, v) \, dv \, dx = \mathbf{m}.$$

Uniform estimates are not directly included in this statement since the function $z \mapsto z \ln(z)$ changes sign. Nevertheless, we can establish such uniform estimates.

COROLLARY 4.2. — *We assume that there exists some (large) $\lambda > 1$ such that*

$$(41) \quad \int \exp(-\lambda \Phi_{\text{ext}}) \, dx < \infty.$$

We suppose also that $0 < \varepsilon \leq 1/(8\lambda)$ and $0 < \theta_\varepsilon \leq 1$. Let $f_\varepsilon^{\text{init}} : \mathbb{R}^N \times \mathbb{R}^N \rightarrow [0, \infty)$ be a sequence of integrable functions that satisfy the following requirements

$$(42) \quad \begin{aligned} & \iint f_\varepsilon^{\text{init}} \, dv \, dx = \mathbf{m}, \\ & \sup_{\substack{0 < \varepsilon \leq 1/(8\lambda) \\ 0 < \theta_\varepsilon \leq 1}} \left\{ \frac{1}{2} \iint |v|^2 f_\varepsilon^{\text{init}} \, dv \, dx + \theta_\varepsilon \iint f_\varepsilon^{\text{init}} |\ln(f_\varepsilon^{\text{init}})| \, dv \, dx \right. \\ & \quad \left. + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon^{\text{init}}|^2 \, dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon^{\text{init}} \, dv \, dx \right\} < \infty, \end{aligned}$$

with

$$\Delta_x \Psi_\varepsilon^{\text{init}} = \frac{1}{\sqrt{\varepsilon}} \left(n_e - \int f_\varepsilon^{\text{init}} \, dv \right).$$

Let $0 < T < \infty$ and let $(f_\varepsilon, \Phi_\varepsilon = \frac{1}{\varepsilon} \Phi_e + \frac{1}{\sqrt{\varepsilon}} \Psi_\varepsilon)$ be the associated solution of (VFP)–(P). Then, uniformly for $0 < \varepsilon \leq 1/(8\lambda)$ and $0 < \theta_\varepsilon \leq 1$:

- (i) $f_\varepsilon(1 + |v|^2 + \theta_\varepsilon |\ln(f_\varepsilon)|) + \varepsilon^{-1} \Phi_e f_\varepsilon$ is bounded in $L^\infty(0, T; L^1(\mathbb{R}^N \times \mathbb{R}^N))$,
- (ii) $\nabla_x \Psi_\varepsilon$ is bounded in $L^\infty(0, T; L^2(\mathbb{R}^N))$,
- (iii) \mathcal{D}_ε is bounded in $L^1(0, T)$.

REMARK 4.3. — In any dimension $N \geq 1$, (41) is always true for quadratic potentials. If $N = 1$ or $N = 2$, hypothesis (41) is satisfied if (h2) is, since $\Phi_{\text{ext}}(x) + \mathbf{m}\Gamma(x) \rightarrow +\infty$ when $|x| \rightarrow +\infty$ and that $\Gamma(x) = -|x|/2$ or $-\ln|x|/(2\pi)$. Therefore, hypothesis (41) needs to be verified only for $N \geq 3$.

Proof. — We first observe that hypothesis (41) implies

$$\int \exp(-\lambda\Phi_e) \, dx < \infty.$$

Indeed, we have $\Phi_e = \Gamma \star n_e - C_* + \Phi_{\text{ext}} \geq \Phi_{\text{ext}} - C_*$. We write, for $h \geq 0$,

$$(43) \quad \begin{aligned} f_\varepsilon \ln(f_\varepsilon) &\leq f_\varepsilon |\ln(f_\varepsilon)| = f_\varepsilon \ln(f_\varepsilon) - 2f_\varepsilon \ln(f_\varepsilon) (\mathbf{1}_{e^{-h} \leq f_\varepsilon \leq 1} + \mathbf{1}_{0 \leq f_\varepsilon < e^{-h}}) \\ &\leq f_\varepsilon \ln(f_\varepsilon) + 2hf_\varepsilon + \frac{4}{e} e^{-h/2}, \end{aligned}$$

and denote

$$E_\varepsilon(f_\varepsilon) = \frac{1}{2} \iint |v|^2 f_\varepsilon \, dv \, dx + \frac{1}{\varepsilon} \iint \Phi_e f_\varepsilon \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx.$$

We now use (43) with $h(x, v) = |v|^2/(8\theta_\varepsilon) + \Phi_e(x)/(4\varepsilon\theta_\varepsilon)$ to infer

$$(44) \quad \begin{aligned} \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) \, dv \, dx &\leq \theta_\varepsilon \iint f_\varepsilon |\ln(f_\varepsilon)| \, dv \, dx \leq \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) \, dv \, dx \\ &\quad + \frac{\theta_\varepsilon}{2} E_\varepsilon(f_\varepsilon) + \theta_\varepsilon \frac{4}{e} \iint \exp(-|v|^2/(16\theta_\varepsilon) - \Phi_e(x)/(8\varepsilon\theta_\varepsilon)) \, dv \, dx. \end{aligned}$$

The last term is equal to

$$\begin{aligned} \theta_\varepsilon \frac{4}{e} \int \exp(-|v|^2/(16\theta_\varepsilon)) \, dv \int \exp(-\Phi_e(x)/(8\varepsilon\theta_\varepsilon)) \, dx \\ \leq \theta_\varepsilon \frac{4}{e} \int \exp(-|v|^2/16) \, dv \int \exp(-\lambda\Phi_e(x)) \, dx, \end{aligned}$$

thus tends to zero as $\theta_\varepsilon \rightarrow 0$ (uniformly for $0 < \varepsilon < 1/(8\lambda)$). Using the dissipation of the entropy given in Proposition 4.1, we then infer

$$\begin{aligned} E_\varepsilon(f_\varepsilon^{\text{init}}) + \theta_\varepsilon \iint f_\varepsilon^{\text{init}} |\ln(f_\varepsilon^{\text{init}})| \, dv \, dx &\geq E_\varepsilon(f_\varepsilon^{\text{init}}) + \theta_\varepsilon \iint f_\varepsilon^{\text{init}} \ln(f_\varepsilon^{\text{init}}) \, dv \, dx \\ &\geq E_\varepsilon(f_\varepsilon) + \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) \, dv \, dx \\ &\geq E_\varepsilon(f_\varepsilon) + \theta_\varepsilon \iint f_\varepsilon |\ln(f_\varepsilon)| \, dv \, dx - \frac{\theta_\varepsilon}{2} E_\varepsilon(f_\varepsilon) + o_{\theta_\varepsilon \rightarrow 0}(1), \end{aligned}$$

and the conclusion follows since $\theta_\varepsilon \leq 1$. □

Until the end of the section, we shall make hypothesis (41). Since we are dealing with the regime

$$0 < \varepsilon \ll 1, \quad 0 < \theta_\varepsilon \ll 1,$$

the estimates in Proposition 4.2 do not provide L^1 -weak compactness on the particle distribution function and its moments; we still need to work with convergences in spaces of finite measures. The first step in the investigation of the asymptotic behavior is summarized in the following claim.

LEMMA 4.4. — We make assumptions (41) and (42). Up to a subsequence, we can assume that f_ε converges to f weakly- \star in $\mathcal{M}^1((0, T) \times \mathbb{R}^N \times \mathbb{R}^N)$. Then, ρ_ε converges to $n_e = \int f \, dv$ in $L^\infty(0, T; H^{-1}(\mathbb{R}^N))$ and in $C^0(0, T; \mathcal{M}^1(\mathbb{R}^N)$ -weak- \star). Moreover, we can assume that

$$J_\varepsilon \rightharpoonup J = \int v f \, dv \quad \text{in } \mathcal{M}^1([0, T] \times \mathbb{R}^N),$$

the limit J is divergence-free and supported in $[0, T] \times \bar{\Omega}$.

Proof. — We follow the arguments of the previous section. We identify the limit of ρ_ε by coming back to the Poisson equation $\sqrt{\varepsilon} \nabla_x \cdot \nabla_x \Psi_\varepsilon = n_e - \rho_\varepsilon$. The time compactness then appears as a consequence of the charge conservation, together with the estimates on the current. We obtain the $L^\infty(0, T; L^1(\mathbb{R}^N))$ estimate on J_ε as in (35). Letting ε go to 0 in the charge conservation equation, we obtain

$$\partial_t n_e + \nabla_x \cdot J = 0 = \nabla_x \cdot J.$$

Still reproducing the arguments of the previous section, based on the conservation of the total charge, we arrive at the following conclusion:

$$\text{Supp}(f) \subset [0, T] \times \bar{\Omega} \times \mathbb{R}^N, \quad \text{Supp}(J) \subset [0, T] \times \bar{\Omega}.$$

Furthermore, J belongs to the set $\mathcal{D}.\mathcal{M}^{\text{ext}}(\mathbb{R}^N)$, it admits a normal trace $J \cdot \nu|_{\partial\Omega}$, which actually vanishes. □

It remains to identify the limit J . As in the case of the pure Vlasov–Poisson equation, the idea consists in introducing a suitable functional intended to compare f_ε to the expected limit. Let $\mathcal{N}_\varepsilon : \mathbb{R}^N \rightarrow (0, \infty)$ be a given function such that

$$\int \mathcal{N}_\varepsilon \, dx = \mathbf{m} = \int n_e \, dx = \iint f(0, x, v) \, dv \, dx$$

and let us set

$$M_{\mathcal{V}, \theta_\varepsilon}(t, x, v) = \frac{1}{(2\pi\theta_\varepsilon)^{N/2}} \exp\left(-\frac{|v - \mathcal{V}(t, x)|^2}{2\theta_\varepsilon}\right).$$

A natural candidate to replace the functional $\mathcal{H}_{\mathcal{V}, \varepsilon}$ would be the relative entropy of f_ε with respect to $n_e(x)M_{\mathcal{V}, \theta_\varepsilon}(t, x, v)$ associated with the non-negative convex function $z \mapsto z \ln(z) - z + 1$, namely

$$\iint \left(f_\varepsilon \ln\left(\frac{f_\varepsilon}{n_e M_{\mathcal{V}, \theta_\varepsilon}}\right) - f_\varepsilon + n_e M_{\mathcal{V}, \theta_\varepsilon} \right) \, dv \, dx,$$

but the first term is clearly meaningless since n_e has compact support. Therefore, we introduce

$$(45) \quad \mathcal{N}_\varepsilon(x) = \frac{\mathbf{m}}{Z_\varepsilon} \exp\left(-\frac{\Phi_e(x)}{\varepsilon\theta_\varepsilon}\right), \quad \text{where } Z_\varepsilon = \int \exp\left(-\frac{\Phi_e(y)}{\varepsilon\theta_\varepsilon}\right) \, dy,$$

and the following modulated functional

$$(46) \quad \mathcal{H}_{\mathcal{V}, \varepsilon}^{\text{FP}} = \theta_\varepsilon \iint \left(f_\varepsilon \ln\left(\frac{f_\varepsilon}{\mathcal{N}_\varepsilon M_{\mathcal{V}, \theta_\varepsilon}}\right) - f_\varepsilon + \mathcal{N}_\varepsilon M_{\mathcal{V}, \theta_\varepsilon} \right) \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx.$$

In fact, $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}$ is up to the term $\frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx$, nothing but the relative entropy of f_ε with respect to $\mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon}$ associated with the non-negative convex function $G : (0, +\infty) \ni z \mapsto z \ln(z) - z + 1$. This implies in particular that the integrand in the first integral of (46) is simply

$$G(f_\varepsilon) - G(\mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon}) - G'(f_\varepsilon)(f_\varepsilon - \mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon}),$$

thus pointwise nonnegative, and vanishes only when $f_\varepsilon = \mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon}$. By definition of \mathcal{N}_ε and $M_{\mathcal{V},\theta_\varepsilon}$ and using the fact

$$\iint f_\varepsilon dv dx = \mathbf{m} = \iint \mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon} dv dx$$

in view of our normalizations, we infer

$$\begin{aligned} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} &= \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) dv dx + \frac{1}{2} \iint |v - \mathcal{V}|^2 f_\varepsilon dv dx + \frac{1}{\varepsilon} \int \Phi_e f_\varepsilon dv dx \\ (47) \quad &+ \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx + \frac{1}{2} N \mathbf{m} \theta_\varepsilon \ln(2\pi\theta_\varepsilon) - \theta_\varepsilon \mathbf{m} \ln(\mathbf{m}/Z_\varepsilon) \\ &= \mathcal{H}_{\mathcal{V},\varepsilon} + \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) dv dx + \frac{1}{2} N \mathbf{m} \theta_\varepsilon \ln(2\pi\theta_\varepsilon) - \theta_\varepsilon \mathbf{m} \ln(\mathbf{m}/Z_\varepsilon). \end{aligned}$$

This second expression of $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}$ justifies the choice we have made for \mathcal{N}_ε . Actually, for our purpose, the exact normalization $\iint \mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon} dv dx = \mathbf{m}$ is not necessary, though natural in a modulated entropy argument, only the fact that the leading order term in $\ln(\mathcal{N}_\varepsilon M_{\mathcal{V},\theta_\varepsilon})$ is $-\Phi_e(x)/(\varepsilon\theta_\varepsilon) - |v - \mathcal{V}|^2/(2\theta_\varepsilon)$ is used. This is related to the fact that the temperature θ_ε is small in the regime we are considering.

Let us now compare $\mathcal{H}_{\mathcal{V},\varepsilon}$ and $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}$ more precisely. As a first step, note that, on the one hand,

$$\theta_\varepsilon \ln(2\pi\theta_\varepsilon) \longrightarrow 0$$

when $\theta_\varepsilon \rightarrow 0$; and on the other hand, that

$$\begin{aligned} |\Omega| &= \int_\Omega \exp\left(-\frac{\Phi_e(y)}{\varepsilon\theta_\varepsilon}\right) dy \leq Z_\varepsilon = \int \exp\left(-\frac{\Phi_e(y)}{\varepsilon\theta_\varepsilon}\right) dy \\ &\leq \int \exp(-\lambda\Phi_e(y)) dy < +\infty \end{aligned}$$

if $\theta_\varepsilon \leq 1$ and $\varepsilon \leq 1/(8\lambda)$, thus, as $\theta_\varepsilon \rightarrow 0$,

$$\theta_\varepsilon \mathbf{m} \ln(\mathbf{m}/Z_\varepsilon) \longrightarrow 0.$$

The inequality (44) implies

$$\begin{aligned} \mathcal{H}_{\mathcal{V},\varepsilon} &= \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} - \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) dv dx - \frac{1}{2} N \mathbf{m} \theta_\varepsilon \ln(2\pi\theta_\varepsilon) + \theta_\varepsilon \mathbf{m} \ln(\mathbf{m}/Z_\varepsilon) \\ (48) \quad &\leq \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} - \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) \mathbf{1}_{f_\varepsilon \leq 1} dv dx + o_{\varepsilon \rightarrow 0}(1) \\ &\leq 2\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + o_{\varepsilon \rightarrow 0}(1). \end{aligned}$$

Then, let us compute the time derivative of the modulated entropy $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}$. We get

$$\begin{aligned} \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} &= \frac{d}{dt} \left\{ \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) dv dx + \frac{1}{2} \iint |v - \mathcal{V}|^2 f_\varepsilon dv dx \right. \\ &\quad \left. + \frac{1}{\varepsilon} \int \Phi_\varepsilon f_\varepsilon dv dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx \right\} \\ &= \frac{d}{dt} \left\{ \theta_\varepsilon \iint f_\varepsilon \ln(f_\varepsilon) dv dx + \frac{1}{2} \iint |v|^2 f_\varepsilon dv dx \right. \\ &\quad \left. + \frac{1}{\varepsilon} \int \Phi_\varepsilon f_\varepsilon dv dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 dx \right\} \\ &\quad + \frac{d}{dt} \left\{ - \iint v \cdot \mathcal{V} f_\varepsilon dv dx + \frac{1}{2} \iint |\mathcal{V}|^2 f_\varepsilon dv dx \right\}. \end{aligned}$$

Bearing in mind the computation for proving Proposition 4.1, we obtain

$$\frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} = -\mathcal{D}_\varepsilon + \frac{d}{dt} \left\{ - \int J_\varepsilon \cdot \mathcal{V} dx + \frac{1}{2} \int \rho_\varepsilon |\mathcal{V}|^2 dx \right\}.$$

Reasoning as in the previous section, and by using the moment equations, we are led to

$$\begin{aligned} \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} &= -\mathcal{D}_\varepsilon + \int \mathcal{V} \cdot J_\varepsilon dx \\ &\quad + \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V}) dx \\ &\quad - \int D\mathcal{V} : (\mathbb{P}_{\mathcal{V},\varepsilon} - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) dx + \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_\varepsilon dx, \end{aligned}$$

using once again that $\nabla_x \cdot (n_\varepsilon \mathcal{V}) = 0$ and the no-flux condition (6). Let us set

$$\mathcal{D}_{\mathcal{V},\varepsilon} = \iint |(v - \mathcal{V}) \sqrt{f_\varepsilon} + 2\theta_\varepsilon \nabla_v \sqrt{f_\varepsilon}|^2 dv dx \geq 0.$$

We rewrite

$$\mathcal{D}_\varepsilon = \mathcal{D}_{\mathcal{V},\varepsilon} - \int \rho_\varepsilon |\mathcal{V}|^2 dx + 2 \int \mathcal{V} \cdot J_\varepsilon dx.$$

Accordingly, we can reorganize terms as follows

$$\begin{aligned} \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} &= -\mathcal{D}_{\mathcal{V},\varepsilon} + \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) \cdot (\partial_t \mathcal{V} + (\mathcal{V} \cdot \nabla_x) \mathcal{V} + \mathcal{V}) dx \\ &\quad - \int D\mathcal{V} : (\mathbb{P}_{\mathcal{V},\varepsilon} - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) dx + \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_\varepsilon dx. \end{aligned}$$

We can summarize the previous manipulations within the following inequality

$$(49) \quad \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + \mathcal{D}_{\mathcal{V},\varepsilon} \leq \frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_\varepsilon dx + r_\varepsilon + \int D\mathcal{V} : (\mathbb{P}_{\mathcal{V},\varepsilon} - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) dx,$$

where, for any $0 < t \leq T$,

$$\int_0^t r_\varepsilon d\tau = \int_0^t \int (\rho_\varepsilon \mathcal{V} - J_\varepsilon) (\partial_\tau \mathcal{V} - \mathcal{V} \cdot \nabla_x \mathcal{V} + \mathcal{V}) dx d\tau$$

tends to 0 as $\varepsilon \rightarrow 0$. We wish to strengthen this result as follows.

LEMMA 4.5. — We make assumptions (41) and (42) and suppose that $\theta_\varepsilon \rightarrow 0$ as $\varepsilon \rightarrow 0$. We have

$$\frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + \mathcal{D}_{\mathcal{V},\varepsilon} \lesssim \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + r_\varepsilon$$

where, for any $0 < t \leq T$, $\lim_{\varepsilon \rightarrow 0} \int_0^t r_\varepsilon \, d\tau = 0$.

Proof. — We can also reproduce the arguments in the previous section used to estimate

$$\frac{1}{\varepsilon} \int \rho_\varepsilon \mathcal{V} \cdot \nabla_x \Phi_\varepsilon \, dx \lesssim \frac{1}{\varepsilon} \int \rho_\varepsilon \Phi_\varepsilon \, dx.$$

For the last term in (49), we have

$$\int D\mathcal{V} : (\mathbb{P}_{\mathcal{V},\varepsilon} - \nabla_x \Psi_\varepsilon \otimes \nabla_x \Psi_\varepsilon) \, dx \leq \|D\mathcal{V}\|_\infty \left(\iint |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx + \int |\nabla_x \Psi_\varepsilon|^2 \, dx \right),$$

so that, using (48),

$$\begin{aligned} \frac{d}{dt} \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + \mathcal{D}_{\mathcal{V},\varepsilon} &\lesssim \frac{1}{\varepsilon} \iint f_\varepsilon \Phi_\varepsilon \, dv \, dx + \frac{1}{2} \iint |v - \mathcal{V}|^2 f_\varepsilon \, dv \, dx + \frac{1}{2} \int |\nabla_x \Psi_\varepsilon|^2 \, dx + r_\varepsilon \\ &\lesssim \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} + r_\varepsilon + o_{\varepsilon \rightarrow 0}(1). \end{aligned}$$

It allows us to conclude by coming back to (49). □

Let us now state our main result concerning the Vlasov-Poisson-Fokker-Planck system. We recall that we may work either with a quadratic potential Φ_{ext} (and then the domain Ω is an ellipsoid), or with a general potential where (h1), (h2), (H1) and (H2) are satisfied.

THEOREM 4.6. — If $N \geq 3$, we make assumption (41), that is we assume that there exists some (large) $\lambda > 1$ such that

$$\int \exp(-\lambda \Phi_{\text{ext}}) \, dx < \infty.$$

Denote by V the solution, on $[0, T]$, to the Lake Equation with friction (LE_f) with the no-flux condition (6) given by Theorem A.1 and consider a smooth extension \mathcal{V} to V . Let

$$f_\varepsilon^{\text{init}} : \mathbb{R}^N \times \mathbb{R}^N \longrightarrow [0, \infty)$$

be a sequence of integrable functions satisfying

$$\iint f_\varepsilon^{\text{init}} \, dv \, dx = \mathbf{m} \quad \text{and} \quad \mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}(f_\varepsilon^{\text{init}}) \longrightarrow 0,$$

where $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}$ is defined in (46). Consider then the associated solutions f_ε of the Vlasov-Poisson-Fokker-Planck equation (VFP)-(P). Then, we have, as $\varepsilon \rightarrow 0$ and $\theta_\varepsilon \rightarrow 0$,

- (i) ρ_ε converges to n_e in $C^0(0, T; \mathcal{M}^1(\mathbb{R}^N)$ -weak- \star);
- (ii) $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}} \rightarrow 0$ uniformly on $[0, T]$;
- (iii) J_ε converges to J in $\mathcal{M}^1([0, T] \times \mathbb{R}^N)$, where $J|_{[0, T] \times \Omega} = V$, $\nabla_x \cdot J = 0$ and $J \cdot \nu(x)|_{\partial\Omega} = 0$.

REMARK 4.7. — We have seen (see Remark 4.3) that the integrability assumption $\int \exp(-\lambda\Phi_{\text{ext}}) dx < \infty$ is automatically satisfied if $N = 1, 2$ by (h2) or for quadratic potentials. When $N \geq 3$, it is also true if Φ_{ext} is convex and tends to $+\infty$ at infinity.

REMARK 4.8. — One may construct an admissible family of initial conditions following the lines of Remark 1.3. In particular, taking G a normalized Gaussian, it is enough to choose θ_ε and σ_ε such that $\theta_\varepsilon \int f_\varepsilon \ln f_\varepsilon \rightarrow 0$, which imposes $\theta_\varepsilon \ln \sigma_\varepsilon \rightarrow 0$.

Proof. — It is clear that if $\mathcal{H}_{\mathcal{V},\varepsilon}^{\text{FP}}(f_\varepsilon^{\text{init}}) \rightarrow 0$, then (42) is satisfied. Item (i) has already been discussed. Applying the Grönwall lemma, we deduce readily that (ii) holds from Lemma 4.5. Coming back to (48), we infer that $\iint |v - \mathcal{V}|^2 f_\varepsilon dv dx$ tends to 0. Then, we appeal to Lemma 3.5 to conclude that J belongs to $L^\infty(0, T; L^2(\mathbb{R}^N))$ and that $J = n_e V$. We can also justify some time-compactness as in the pure Vlasov–Poisson case. \square

APPENDIX A. SMOOTH SOLUTIONS OF THE LAKE EQUATIONS

THEOREM A.1. — Let Ω be a smooth ($\partial\Omega$ of class C^{s+1} is enough) bounded open set in \mathbb{R}^N , γ be a real constant, $s \in \mathbb{N}$ such that $s > 1 + N/2$ and $n_e : \Omega \rightarrow \mathbb{R}$ in H^{s+1} such that $\inf_\Omega n_e > 0$. Let $V^{\text{init}} : \Omega \rightarrow \mathbb{R}^N$ be a divergence free vector field in H^s satisfying the no-flux condition $V^{\text{init}} \cdot \nu = 0$ on $\partial\Omega$. There exists $T > 0$ and a unique solution $V \in L^\infty(0, T; H^s(B(0, R)))$ of

$$\begin{cases} \partial_t V + V \cdot \nabla_x V + \nabla_x p = -\gamma V, \\ \nabla_x \cdot (n_e V) = 0, \end{cases}$$

with the no-flux condition (6). Moreover, we have

$$\sup_{0 \leq t \leq T} \left(\|V(t)\|_{H^s} + \|\partial_t V(t)\|_{H^{s-1}} + \|\nabla_x p(t)\|_{H^s} + \|\partial_t \nabla_x p(t)\|_{H^{s-1}} \right) \leq C(T)$$

for some positive constant $C(T)$ depending on γ, T, n_e and the initial datum.

Proof. — The scheme of proof is exactly the same as in [44]. We shall denote $\widehat{V} := n_e V$, which is divergence free. Applying $\nabla \cdot (n_e \cdot)$ to the equation, we see that the pressure p satisfies, for any t , the elliptic equation

$$\begin{aligned} -\nabla_x \cdot (n_e \nabla_x p) &= \nabla_x \cdot (n_e V \cdot \nabla_x V) = \nabla_x \cdot \left(\widehat{V} \cdot \nabla_x \left(\frac{1}{n_e} \widehat{V} \right) \right) \\ (50) \quad &= \widehat{V} \cdot \nabla_x \left(\nabla_x \cdot \left(\frac{1}{n_e} \widehat{V} \right) \right) + \sum_{1 \leq j, k \leq N} \partial_j \widehat{V}_k \partial_k (\widehat{V}_j / n_e) \\ &= \widehat{V} \cdot \nabla_x \left(\widehat{V} \cdot \nabla_x (1/n_e) \right) + \sum_{1 \leq j, k \leq N} \partial_j \widehat{V}_k \partial_k (\widehat{V}_j / n_e), \end{aligned}$$

where we have used that \widehat{V} is divergence free as well as the identity

$$\nabla_x \cdot (\mathcal{V} \cdot \nabla_x \mathcal{U}) - \mathcal{V} \cdot \nabla_x (\nabla_x \cdot \mathcal{U}) = \sum_{1 \leq j, k \leq N} \partial_j \mathcal{V}_k \partial_k \mathcal{U}_j.$$

We may further impose a suitable Neumann boundary condition for p on $\partial\Omega$. We recall that for $\sigma > N/2$, H^σ is an algebra. Notice that if $V \in H^s$, with $s > 1 + N/2$, then the right-hand side of (50) is in H^{s-1} and

$$\left\| \widehat{V} \cdot \nabla_x \left(\widehat{V} \cdot \nabla_x (1/n_e) \right) + \sum_{1 \leq j, k \leq N} \partial_j \widehat{V}_k \partial_k \left(\widehat{V}_j / n_e \right) \right\|_{H^{s-1}} \leq C \|V\|_{H^s}^2,$$

where C depends on $\inf_\Omega n_e$ (which is assumed positive) and the H^{s+1} norm of n_e .

Since n_e is in H^{s+1} and bounded away from zero and since the boundary is assumed of class C^{s+1} , it follows from classical elliptic estimates that (50) endowed with the Neumann condition on $\partial\Omega$ has a solution $p \in H^{s+1}(\Omega)$, defined up to a constant in each connected component of Ω , enjoying the estimate

$$(51) \quad \|\nabla p\|_{H^s} \leq C \|V\|_{H^s}^2,$$

where C depends on $\inf_\Omega n_e$ and the H^{s+1} norm of n_e . Assume now that V is a smooth solution of (A.1) and let us perform an H^s estimate. For any $\alpha \in (\mathbb{N} \cup \{0\})^d$ with $|\alpha| \leq s$, we have

$$\begin{aligned} & \frac{d}{dt} \int_\Omega |\partial^\alpha V|^2 dx \\ &= -2 \int_\Omega \partial^\alpha V \cdot \partial^\alpha ((V \cdot \nabla_x) V) dx - 2 \int_\Omega \partial^\alpha V \cdot \partial^\alpha \nabla_x p dx - 2\gamma \int_\Omega |\partial^\alpha V|^2 dx. \end{aligned}$$

Using classical commutator estimates, the Sobolev imbedding $H^s \subset W^{1,\infty}$ and the H^s estimate (51) on ∇p , we then deduce

$$\frac{d}{dt} \int_\Omega |\partial^\alpha V|^2 dx \leq -2 \int_\Omega \partial^\alpha V \cdot ((V \cdot \nabla_x) \partial^\alpha V) dx + C(\|V\|_{H^s} + \|V\|_{H^s}^2 + \|V\|_{H^s}^3).$$

We use integration by parts for the first integral (recall that $V \cdot \nu = 0$ on the boundary), which then becomes $\int_\Omega (\nabla_x \cdot V) |\partial^\alpha V|^2 dx \leq C \|V\|_{H^s}^3$. This yields

$$\frac{d}{dt} \int_\Omega |\partial^\alpha V|^2 dx \leq C(\|V\|_{H^s} + \|V\|_{H^s}^2 + \|V\|_{H^s}^3),$$

and it follows that, for some $T_0 > 0$ depending only on γ , n_e and V^{init} , we have $\|V\|_{L^\infty(0, T_0; H^s)} \leq 2\|V^{\text{init}}\|_{H^s}$. The conclusion of the theorem follows from a suitable viscous approximation where a careful treatment of boundary terms is needed, see [45]. \square

APPENDIX B. CONSTRUCTION OF AN EXTENDED DIVERGENCE–FREE VELOCITY

LEMMA B.1. — *Let $V \in L^\infty(0, T; H^s(B(0, R), \mathbb{R}^N))$ be a divergence free vector field in H^s , with $s > 1 + N/2$, satisfying the no flux condition $V \cdot \nu = 0$ on $\partial B(0, R)$. There exists a solenoidal extension \mathcal{V} of the vector field V defined on the whole space and compactly supported. Namely, $\mathcal{V} \in L^\infty(0, T; H^s(\mathbb{R}^N, \mathbb{R}^N))$ and it satisfies:*

- (i) (7),
- (ii) $\nabla_x \cdot \mathcal{V} = 0$ in \mathbb{R}^N .

Proof. — Let us assume $N = 2$ or $N = 3$. Since $\nabla_x \cdot V = 0$ in the ball $B(0, R)$, which is convex, there exists $h \in L^\infty(0, T; H^{s+1}(B(0, R), \mathbb{R}^N))$ such that $V = \nabla \times h$. Then, by standard extension results (see e.g. [33, Chap. I: Th. 2.1 p. 17 & Th. 8.1 p. 42]), there exists an extension $\tilde{h} \in L^\infty(0, T; H^{s+1}(\mathbb{R}^N, \mathbb{R}^N))$ to h . Considering a cut-off function $\chi \in C_c^\infty(\mathbb{R}^N)$ such that $\chi(x) = 1$ for $x \in B(0, 3R/2)$ and denoting $\mathcal{V} = \nabla \times (\chi \tilde{h})$, we see that \mathcal{V} enjoys the desired properties. For $N \geq 4$, the construction is similar but involves differential forms. \square

The arguments generalize readily to the case where the domain is an ellipsoid.

REMARK B.2. — In the case of a general potential, n_e is not uniform, and it is possible to construct an extension $\mathcal{V} \in L^\infty(0, T; H^s(\mathbb{R}^N, \mathbb{R}^N))$ which satisfies:

- (i) (7),
- (ii) $\nabla_x \cdot (n_e \mathcal{V}) = 0$ in \mathbb{R}^N .

However, it requires further topological hypotheses on Ω . Assuming (H1), and assuming also that \mathcal{K} is connected, and $\partial\Omega$ has a finite number of connected components, we may apply [27, Co. 3.2]: $n_e V$ is divergence free in the smooth domain Ω , hence we can construct a divergence free extension

$$\mathcal{J} : [0, T] \times \mathbb{R}^N \longrightarrow \mathbb{R}^N$$

to $n_e V$. Using a cut-off function, we may take \mathcal{J} compactly supported in an arbitrary neighborhood of \mathcal{K} , the latter can be chosen so that $\Delta\Phi_{\text{ext}}$ remains > 0 . Finally, we set $\mathcal{V} = \mathcal{J}/(\Delta\Phi_{\text{ext}})$, which is well-defined even when $\Delta\Phi_{\text{ext}}$ vanishes.

APPENDIX C. ABOUT THE EXPERIMENTAL PARAMETERS IN A MAGNETO-OPTICAL TRAP EXPERIMENT

The goal of this section is to present an experimental situation where the asymptotic limit studied in this article is relevant.

Atoms in a magneto-optical trap (MOT) interact in a complicated way with lasers, whose frequency is close to an atomic resonance. It is customary to split this complicated interaction into several parts: a friction, which is linear in velocity as a first approximation; an anisotropic external trap, which is harmonic as a first approximation; a velocity diffusion; a repulsive interaction [46], usually modeled as Coulombic; an attractive interaction of a rather complicated form [10]. The latter interaction is always weaker than the repulsive one, so that we neglect it here. Putting all these effects together yields the Vlasov-Fokker-Planck equation for $f(t, x, v)$, the density of particles in phase space:

$$(52) \quad \partial_t f + v \cdot \nabla_x f - \nabla_x \phi \cdot \nabla_v f = \nabla_v \cdot (\kappa v f + D \nabla_v f),$$

$$(53) \quad \Delta_x \phi = 3\omega_0^2 - c\rho,$$

where ω_0 is the trap frequency, κ the friction, D the velocity diffusion coefficient, and c the "pseudo-Coulombic" coupling constant. Here the integral of f is normalized

to 1 and the dimension of f is $L^{-2N}T^{-1}$ (L stands for "length" and T for "time"). For the sake of concreteness we assume $N = 3$. Since our purpose is only to discuss orders of magnitudes here, we have taken an isotropic trap. This model is clearly very approximate, but accounts for some basic experimental results, and plays the role of a reference model to which more refined measurements and theoretical descriptions are compared (see for instance [7] for a recent comparison between models and theory in very large MOTs).

We write dimensionless equations using $t = T\tilde{t}$, $x = L\tilde{x}$, $f = L^{-6}T^3\tilde{f}$ and $\phi = L^2T^{-2}\tilde{\phi}$; we then drop the $\tilde{}$ for convenience. We obtain

$$(54) \quad \partial_t f + v \cdot \nabla_x f - \nabla_x \phi \cdot \nabla_v f = \kappa T \nabla_v \cdot \left(v f + \frac{DT^2}{\kappa L^2} \nabla_v f \right),$$

$$(55) \quad \Delta_x \phi = 3\omega_0^2 T^2 \left(1 - \frac{c}{3\omega_0^2 L^3} \rho \right).$$

We now choose

$$L = \left(\frac{c}{3\omega_0^2} \right)^{1/3}.$$

Equation (55) becomes

$$\Delta_x \phi = 3\omega_0^2 T^2 (1 - \rho).$$

Hence in the strong coupling limit, $\rho = 1$, and L is essentially the size of the atomic cloud. We now introduce

$$L_g = \left(\frac{D}{\kappa \omega_0^2} \right)^{1/2} \quad \text{and} \quad Q = \frac{\omega_0}{\kappa}.$$

The length L_g would be the size of the cloud for a vanishingly weak repulsion; in this case, the density profile is a Gaussian. The parameter Q is the quality factor of the damped harmonic oscillations in the trap. We also introduce

$$\varepsilon = \frac{1}{3} \left(\frac{L_g}{L} \right)$$

and choose

$$T = \frac{1}{\omega_0} \left(\frac{L}{L_g} \right)^{1/2}.$$

The system (54)-(55) becomes

$$(56) \quad \partial_t f + v \cdot \nabla_x f - \nabla_x \phi \cdot \nabla_v f = \kappa T \nabla_v \cdot \left(v f + \frac{L_g}{L} \nabla_v f \right),$$

$$(57) \quad \Delta_x \phi = \frac{1}{\varepsilon} (1 - \rho).$$

This corresponds to the system (VFP) we have studied, with an dimensionless friction coefficient $\gamma = \kappa T$, instead of $\gamma = 1$. A typical value for L_g is $200 \mu m$, and Q is at most of order $1^{(2)}$. The size of the cloud L varies depending on the number of trapped

⁽²⁾We thank R. Kaiser and G. Labeyrie for providing us these informations.

atoms, a typical value for the regime we would like to describe being $L \sim 2 \text{ mm}$ (see [7, Fig. 5]). Putting these orders of magnitude together, we obtain:

$$\varepsilon = \frac{1}{30}, \quad \theta = \frac{1}{10}, \quad \gamma \simeq 3.$$

It is thus not unreasonable to describe such a system by the limit regime studied in this paper: $\varepsilon \rightarrow 0$, $\theta \rightarrow 0$. The typical reduced friction in a MOT experiment would however likely be larger than 1.

REFERENCES

- [1] A. A. ARSENEV – “Global existence of a weak solution of Vlasov’s system of equations”, *U.S.S.R. Comput. Math. and Math. Phys.* **15** (1975), p. 131–143.
- [2] G. BEN AROUS & A. GUIONNET – “Large deviations for Wigner’s law and Voiculescu’s non-commutative entropy”, *Probab. Theory Relat. Fields* **108** (1997), no. 4, p. 517–542.
- [3] F. BOUCHUT – “Existence and uniqueness of a global smooth solution for the VPFP system in three dimensions”, *J. Funct. Anal.* **111** (1993), p. 239–258.
- [4] Y. BRENIER – “A homogenized model for vortex sheets”, *Arch. Rational Mech. Anal.* **138** (1997), no. 4, p. 319–353.
- [5] ———, “Convergence of the Vlasov-Poisson system to the incompressible Euler equations”, *Comm. Partial Differential Equations* **25** (2000), no. 3-4, p. 737–754.
- [6] L. CAFFARELLI – “The obstacle problem revisited”, *J. Fourier Anal. Appl.* **4** (1998), no. 3&4, p. 383–402.
- [7] A. CAMARA, R. KAISER & G. LABEYRIE – “Behavior of a very large magneto-optical trap”, *Phys. Rev. A* **90** (2014), no. 6, p. 063404.
- [8] D. CHAFAÏ, N. GOZLAN & P.-A. ZITT – “First order asymptotics for confined particles with singular pair repulsions”, *Ann. Appl. Probab.* **24** (2014), no. 6, p. 2371–2413.
- [9] G.-Q. CHEN & H. FRID – “Extended divergence-measure fields and the Euler equations for gas dynamics”, *Comm. Math. Phys.* **236** (2003), p. 251–280.
- [10] J. DALIBARD – “Laser cooling of an optically thick gas: The simplest radiation pressure trap?”, *Optics Comm.* **68** (1988), p. 203–208.
- [11] J. DALIBARD & C. COHEN-TANNOUDJI – “Atomic motion in laser light: connection between semiclassical and quantum descriptions”, *J. Phys. B* **18** (1985), no. 8, p. 1661.
- [12] P. DEGOND – “Global existence of smooth solutions for the Vlasov-Fokker-Planck equation in 1 and 2 space dimension”, *Ann. Sci. École Norm. Sup. (4)* **19** (1986), p. 519–542.
- [13] J. DOLBEAULT – “Free energy and solutions of the Vlasov-Poisson-Fokker-Planck system: external potential and confinement (large time behavior and steady states)”, *J. Math. Pures Appl.* **78** (1999), no. 2, p. 121–157.
- [14] D. DUBIN & T. O’NEIL – “Trapped nonneutral plasmas, liquids, and crystals (the thermal equilibrium states)”, *Rev. Modern Phys.* **71** (1999), no. 1, p. 87–172.
- [15] J. FREHSE – “On the regularity of the solution of a second order variational inequality”, *Boll. Un. Mat. Ital. (4)* **6** (1972), p. 312–315.
- [16] A. FRIEDMAN – *Variational principles and free-boundary problems*, Pure and Applied Math., John Wiley & Sons, Inc., New York, 1982, A Wiley-Interscience Publication.
- [17] A. FRIEDMAN & D. PHILLIPS – “The free boundary of a semilinear elliptic equation”, *Trans. Amer. Math. Soc.* **282** (1984), no. 1, p. 153–182.
- [18] O. FROSTMAN – “Potentiel d’équilibre et capacité des ensembles avec quelques applications à la théorie des fonctions”, *Meddelanden Mat. Sem. Univ. Lund* **3** (1935), p. 1–118.
- [19] M. A. FURMAN – “Compact complex expressions for the electric field of two-dimensional elliptical charge distributions”, *Amer. J. Phys.* **62** (12) (1994), p. 1134–1140.
- [20] R. T. GLASSEY – *The Cauchy problem in kinetic theory*, Society for Industrial and Applied Mathematics (SIAM), Philadelphia, PA, 1996.

- [21] E. GRENIER – “Defect measures of the Vlasov–Poisson system in the quasineutral regime”, *Comm. Partial Differential Equations* **20** (1995), no. 7-8, p. 1189–1215.
- [22] ———, “Oscillations in quasineutral plasmas”, *Comm. Partial Differential Equations* **21** (1996), no. 3-4, p. 363–394.
- [23] D. HAN-KWAN – “Quasineutral limit of the Vlasov–Poisson system with massless electrons”, *Comm. Partial Differential Equations* **36** (2011), no. 8, p. 1385–1425.
- [24] D. HAN-KWAN & M. HAURAY – “Stability issues in the quasineutral limit of the one-dimensional Vlasov–Poisson equation”, *Comm. Math. Phys.* **334** (2015), no. 2, p. 1101–1152.
- [25] D. HAN-KWAN & M. IACOBELLI – “Quasineutral limit for Vlasov–Poisson via Wasserstein stability estimates in higher dimension”, Preprint École polytechnique, [arXiv:1503.06097](https://arxiv.org/abs/1503.06097), 2015.
- [26] H. HEDENMÅLM & N. MAKAROV – “Coulomb gas ensembles and Laplacian growth”, *Proc. London Math. Soc. (3)* **106** (2013), no. 4, p. 859–907.
- [27] T. KATO, M. MITREA, G. PONCE & M. TAYLOR – “Extension and representation of divergence-free vector fields”, *Math. Res. Lett.* **7** (2000), p. 643–650.
- [28] O. D. KELLOGG – *Foundations of potential theory*, Grundlehren Math. Wiss., vol. 31, Springer-Verlag, Berlin, New York, 1967, Reprint from the first edition of 1929.
- [29] D. KINDERLEHRER & L. NIRENBERG – “Regularity in free boundary problems”, *Ann. Scuola Norm. Sup. Pisa Cl. Sci. (4)* **4** (1977), no. 2, p. 373–391.
- [30] D. KINDERLEHRER & G. STAMPACCHIA – *An introduction to variational inequalities and their applications*, Pure and Applied Math., vol. 88, Academic Press, Inc., New York, London, 1980.
- [31] C. D. LEVERMORE, M. OLIVER & E. S. TITI – “Global well-posedness for models of shallow water in a basin with a varying bottom”, *Indiana Univ. Math. J.* **45** (1996), no. 2, p. 479–510.
- [32] ———, “Global well-posedness for the lake equations”, *Phys. D* **98** (1996), no. 2-4, p. 492–509.
- [33] J.-L. LIONS & E. MAGENES – *Problèmes aux limites non homogènes et applications (volume 1)*, Travaux et recherches mathématiques, Dunod, Paris, 1968.
- [34] P.-L. LIONS & B. PERTHAME – “Propagation of moments and regularity for the 3-dimensional Vlasov–Poisson system”, *Invent. Math.* **105** (1991), p. 415–430.
- [35] N. MASMOUDI – “From Vlasov–Poisson system to the incompressible Euler system”, *Comm. Partial Differential Equations* **26** (2001), no. 9-10, p. 1913–1928.
- [36] ———, “Rigorous derivation of the anelastic approximation”, *J. Math. Pures Appl.* **88** (2007), no. 3, p. 230–240.
- [37] J. MENDONÇA, R. KAISER, H. TERÇAS & J. LOUREIRO – “Collective oscillations in ultra-cold atomic gas”, *Phys. Rev. A* **78** (2008), no. 1, p. 013408.
- [38] Y. OGURA & N. PHILLIPS – “Scale analysis for deep and shallow convection in the atmosphere”, *J. Atmospheric Sci.* **19** (1962), p. 173–179.
- [39] M. OLIVER – “Classical solutions for a generalized Euler equation in two dimensions”, *J. Math. Anal. Appl.* **215** (1997), no. 2, p. 471–484.
- [40] A. OLIVETTI – “Effets collectifs et particules en interaction : des systèmes à longue portée aux atomes froids”, PhD Thesis, Univ. Nice Sophia Antipolis, 2011.
- [41] K. PEAFFELMOSER – “Global existence of the Vlasov–Poisson system in three dimensions for general initial data”, *J. Differential Equations* **95** (1992), p. 281–303.
- [42] E. SAFF & V. TOTIK – *Logarithmic potentials with external fields*, Grundlehren Math. Wiss., Springer, Berlin, 1997.
- [43] S. SERFATY – “Coulomb gases and Ginzburg–Landau vortices”, 2014, Lecture Notes of the “nachdiplom vorlesung” course, ETH Zürich, Forschungsinstitut für Mathematik.
- [44] R. TEMAM – “On the Euler equations of incompressible perfect fluids”, *J. Functional Analysis* **20** (1975), no. 1, p. 32–43.
- [45] ———, “Local existence of C^∞ solutions of the Euler equations of incompressible perfect fluids”, in *Turbulence and Navier–Stokes equations (Orsay, 1975)*, Lect. Notes in Math., vol. 565, Springer, Berlin, 1976, p. 184–194.
- [46] T. WALKER, D. SESKO & C. WIEMAN – “Collective behavior of optically trapped neutral atoms”, *Phys. Rev. Lett.* **64** (1990), p. 408–412.
- [47] D. WINELAND, J. BOLLINGER, W. ITANO & J. PRESTAGE – “Angular momentum of trapped atomic particles”, *J. Opt. Soc. Amer.* **2** (1985), no. 11, p. 1721–1730.

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