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The three-dimensional finite Larmor radius approximation

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Abstract

Following Frénod and Sonnendrücker ([12]), we consider the finite Larmor radius regime for a plasma submitted to a large magnetic field and take into account both the quasineutrality and the local thermodynamic equilibrium of the electrons. We then rigorously establish the asymptotic gyrokinetic limit of the rescaled and modified Vlasov-Poisson system in a three-dimensional setting with the help of an averaging lemma.

Keywords: Gyrokinetic approximation - Vlasov-Poisson equation - Finite Larmor Radius scaling - Averaging lemma.

1 Introduction and main results

1.1 Physical motivation

We are interested in the behaviour of a plasma (*id est* a gaz made of ions with individual charge Ze and mass m_i and electrons with individual charge $-e$ and mass m_e , with $m_i \gg m_e$) which is submitted to a large external magnetic field. It is “well-known” that such a field induces fast small oscillations for the particles and consequently introduces a new small time scale which is very restrictive and inconvenient from the numerical point of view. The simulation of such plasmas appears to be primordial since the model can be applied to tokamak plasmas from magnetic confinement fusion (like for the ITER project).

1.1.1 Heuristic study

Let us give some heuristic formal arguments to investigate the behaviour of the plasma: if we consider the motion of one particle (of charge $q > 0$, mass m , position x and velocity v) submitted to an external constant field B , the fundamental principle of mechanics gives that:

$$\frac{dx}{dt} = v, \quad \frac{dv}{dt} = \frac{q}{m}(v \wedge B) \quad (1.1)$$

Straightforward calculations show first of all that the parallel velocity, denoted by v_{\parallel} (that is to say the component of the velocity in the direction of the magnetic field) is conserved and thanks to the conservation of the kinetic energy, so is the norm of the perpendicular velocity v_{\perp} (the component of the velocity in the perpendicular plane). Actually, we can see that the particle moves on a helix whose axis is the direction of the magnetic field. The rotation period (around the axis) is the inverse of the cyclotron frequency Ω :

$$\Omega = \frac{|q||B|}{m} \quad (1.2)$$

and the radius is the so-called Larmor radius:

$$r_L = \frac{|v_{\perp}|}{\Omega} \quad (1.3)$$

In the case where the magnetic field is very strong, Ω tends to infinity whereas r_L tends to zero. More precisely, if we take $|B| \sim \frac{1}{\epsilon}$ (with $\epsilon \rightarrow 0$) we have:

$$\begin{cases} \Omega \sim \frac{1}{\epsilon} \\ r_L \sim \epsilon \end{cases}$$

The approximation which consists in considering $r_L = 0$ is the classical guiding center approximation ([17]). This means that each particle is assimilated to its “guiding center” (in other words its “instantaneous rotation center”), which is equivalent to neglect the very fast rotation of the particle around the axis.

If one also applies some external constant electric field E , a similar computation shows that there appears:

1. an acceleration $\frac{E \cdot B}{|B|}$ in the direction of B . If we consider $E \sim 1$, then:

$$\frac{E \cdot B}{|B|} \sim 1 \quad (1.4)$$

2. a drift $\frac{E \wedge B}{|B|^2}$ in the orthogonal plane. We have:

$$\frac{E \wedge B}{|B|^2} \sim \epsilon \quad (1.5)$$

This drift, usually called the electric drift is problematic as regards to the issue of plasma confinement. It is negligible compared to the acceleration in the direction of B , but in the time scale for plasma fusion which is expected to be very long, one can not neglect this small drift, since it creates a displacement of order ϵt (t represents the time).

At last, note also that if the fields are not constant, various other drifts may appear, whose order in ϵ is higher than those of the electric drift.

Actually, the fields considered are neither constant, nor external, but self-induced by the plasma itself. The effects we would like to describe are due to the non-linear interaction between the particles and the electromagnetic field.

1.1.2 The mathematical model

In all the sequel, we assume that the magnetic field is external and constant and we suppose that the speed of particles is small compared to the speed of light, so that we can use the electrostatic approximation which consists in reducing the Maxwell equations

to the Poisson equation. Finally, we decide to opt for a kinetic description for the ions: in other words, the time and space scales considered here are such that ions are not at a thermodynamic equilibrium and their density is governed by a kinetic equation.

The basic model usually considered for the ions is the following Vlasov-Poisson system:

$$\begin{cases} \partial_t f + v \cdot \nabla_x f + (E + v \wedge B) \cdot \nabla_v f = 0 \\ E = -\nabla_x V \\ -\Delta_x V = \int f dv \\ f_{t=0} = f_0 \end{cases}$$

where $f(t, x, v)$ is the density of ions, with $t \in \mathbb{R}^+$, $x \in \mathbb{R}^d$ or $\mathbb{R}^d/\mathbb{Z}^d$, $v \in \mathbb{R}^d$ (usually $d = 2$ or 3), meaning that $f(t, x, v)dx dv$ gives the number of ions in the infinitesimal volume $[x, x + dx] \times [v, v + dv]$ at time t (note that in this model, electrons are for the moment neglected).

1.1.3 The gyrokinetic approximation

It is important from a numerical point of view to establish the asymptotic equation when $|B|$ tends to infinity. Indeed, we expect the asymptotic equation to be “easier” to handle: only one time and space scale, perhaps less variables in the phase space to deal with... The derivation of such equations is usually referred to in the mathematic literature as “gyrokinetic approximation”.

Rigorous justifications of these derivations with various time and space observation scales have only appeared at the end of the nineties. We refer for instance to the works of Brenier ([5]), Frénod and Sonnendrücker ([11]-[12]), Frénod, Raviart and Sonnendrücker ([10]), Golse and Saint-Raymond ([15]-[16]), Saint-Raymond ([21]-[22]).

The classical “guiding center approximation” corresponds to the following scaling for the Vlasov-Poisson system (from now on and until the end of the paper, B is a constant vector, say for instance $B = \frac{1}{\epsilon} e_z$):

$$\begin{cases} \partial_t f_\epsilon + v \cdot \nabla_x f_\epsilon + (E_\epsilon + \frac{v \wedge e_z}{\epsilon}) \cdot \nabla_v f_\epsilon = 0 \\ E_\epsilon = -\nabla_x V_\epsilon \\ -\Delta_x V_\epsilon = \int f_\epsilon dv \\ f_{\epsilon, t=0} = f_0 \end{cases} \quad (1.6)$$

The articles [11] and [15] show that when $\epsilon \rightarrow 0$, this leads to a one-dimensional kinetic equation in the direction of B :

$$\begin{cases} \partial_t f + v_\parallel \cdot \nabla_x f + E_\parallel \cdot \nabla_v f = 0 \\ E = -\nabla_x V \\ -\Delta_x V = \int f dv \\ f_{t=0} = f_0 \end{cases} \quad (1.7)$$

Notice that the electric drift does not appear; this was expected since we have seen in the formal analysis that this drift was of higher order in ϵ than the other effects. This shows in particular that this approximation is not sufficient for the numerical simulation of tokamaks. In order to make this drift appear, there exists to our knowledge two main possibilities:

1. one consists in restricting to a $2D$ problem in the plane orthogonal to B ([15]),
2. the other consists in rescaling the orthogonal scales in order to get both transport and electric drift at the same order ([12]).

This work directly follows the articles [12] and [10] where the authors considered the “finite Larmor radius approximation”. This means that the spatial observation scale in the plane orthogonal to B is chosen smaller than the one in the parallel direction, more precisely with the same order as the Larmor radius r_L , so that one can expect the electric drift to appear in the asymptotic equation.

In some sense, having such a scaling allows the electric field to significantly vary across a Larmor radius, which is not the case for instance in (1.6). Moreover, in this situation, the positions of the particles are no longer assimilated to the position of their “guiding center” and we will have to perform an average over one fast oscillation period (the so-called gyroaverage) in order to get a sort of averaged number density.

1.2 Scaling and existing results

The system we are going to study is based on the “finite Larmor radius scaling” and takes into account the quasineutrality of the plasma.

1.2.1 The (refined) mathematical model

We refer to [12] for a complete discussion on the scaling. Let us recall briefly and quite crudely how it works.

Let L_{\parallel} be the characteristic length in the direction of the magnetic field and L_{\perp} be the characteristic length in the perpendicular plane. We consider that $L_{\parallel} \sim 1$ and $L_{\perp} \sim \epsilon$ and define the dimensionless variables $x'_{\parallel} = \frac{x_{\parallel}}{L_{\parallel}}$ and $x'_{\perp} = \frac{x_{\perp}}{L_{\perp}}$. In the same fashion we also define the dimensionless variables t' and v' with characteristic time and velocities with the same order as L_{\parallel} and introduce the new number density f' defined by $\bar{f}f'(t', x', v') = f(t, x, v)$ (and we define likewise the new electric field and potential $\bar{E}E'(t', x', v') = E(t, x, v)$ and $\bar{V}V'(t', x', v') = V(t, x, v)$). We consider the scaling $\bar{f}, \bar{E} \sim 1$ and $\bar{V} \sim \epsilon$. At last, we introduce the Debye length of the plasma λ_D , which appears in the Poisson equation. In order to take into account the quasineutrality of the plasma, we take from now on $\lambda_D \sim \sqrt{\epsilon}$.

The Poisson equation states in this scaling:

$$-\epsilon \Delta_{x'_{\parallel}} V'_{\epsilon} - \frac{1}{\epsilon} \Delta_{x'_{\perp}} V'_{\epsilon} = \frac{1}{\epsilon} (n_{\epsilon}^i - n_{\epsilon}^e) \quad (1.8)$$

where $n_{\epsilon}^i = \int f'_{\epsilon} dv'$ is the density of ions and n_{ϵ}^e the density of electrons. The density distribution of ions is normalized so that $\int f'_0 dv' dx' = 1$.

The main difference between Frénod and Sonnendrücker’s model and ours lies in the following. Instead of considering a fixed background of electrons, and since $\frac{m_e}{m_i} \ll 1$, we make the usual assumption that the (adiabatic) electrons are instantaneously at a local thermodynamic equilibrium, so that their density follows a Boltzmann-Maxwell distribution:

$$n_{\epsilon}^e(x, t) = \exp\left(\frac{eV'_{\epsilon}}{k_B T_e}\right) \quad (1.9)$$

where k_B is the Boltzmann constant, $-e$ the charge and T_e the temperature of the electrons. We consider that $\frac{e}{k_B T_e} \sim 1$.

We make the assumption that we are not far from a fixed background of electrons, so that we can linearize this expression:

$$n_{\epsilon}^e(x, t) = 1 + V'_{\epsilon} \quad (1.10)$$

We are obviously aware that this assumption is not really satisfactory both from a mathematical and physical point of view; let us just consider this as a plain mathematical model which allows us to give rigorous justifications in this work. The problem of a fixed background of electrons, i.e. $n_\epsilon^e = 1$, brings actually more interesting formal results; this point will be discussed in the last section.

The Poisson equation can now be written:

$$V'_\epsilon - \epsilon^2 \Delta_{x'_\parallel} V'_\epsilon - \Delta_{x'_\perp} V'_\epsilon = \int f_\epsilon dv' - \int f_0 dv' dx' \quad (1.11)$$

The dimensionless system (1.6) becomes (for the sake of simplicity, we forget the primes):

$$\begin{cases} \partial_t f_\epsilon + \frac{v_\perp}{\epsilon} \cdot \nabla_x f_\epsilon + v_\parallel \cdot \nabla_x f_\epsilon + (E_\epsilon + \frac{v \wedge B}{\epsilon}) \cdot \nabla_v f_\epsilon = 0 \\ E_\epsilon = (-\nabla_{x_\perp} V_\epsilon, -\epsilon \nabla_{x_\parallel} V_\epsilon) \\ V_\epsilon - \epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon = \int f_\epsilon dv - \int f_0 dv dx \\ f_{\epsilon, t=0} = f_{\epsilon, 0} \end{cases} \quad (1.12)$$

with the notation $\Delta_{x_\parallel} = \partial_{x_\parallel}^2$ and $\Delta_{x_\perp} = \Delta - \Delta_{x_\parallel}$,

the problem being posed for $(x_\perp, x_\parallel, v) \in \mathbb{T}^2 \times \mathbb{T} \times \mathbb{R}^3$ (with $\mathbb{T} = \mathbb{R}/\mathbb{Z}$ equipped with the restriction of the Lebesgue measure to $[0, 1]$).

1.2.2 State of the art about the Finite Larmor Radius Approximation

Using homogenization arguments, Frénod and Sonnendrücker established the convergence in some weak sense of sequences of solutions $(f_\epsilon)_{\epsilon \geq 0}$ of similar systems, in two cases, namely in some pseudo 2D case (assuming that nothing depends on x_\parallel and v_\parallel) and in a 3D case when the electric field is external. The main tool used to establish the convergence is the “2-scale convergence” introduced by Nguetseng [20] and Allaire [2] that we will recall later on.

1. The 3D case:

Assume that we deal with an external electric field $E_\epsilon = E \in \mathcal{C}^1(\mathbb{R} \times \mathbb{R}^3)$:

$$\begin{cases} \partial_t f_\epsilon + v_\parallel \cdot \nabla_x f_\epsilon + \frac{v_\perp}{\epsilon} \cdot \nabla_x f_\epsilon + (E + \frac{v \wedge e_z}{\epsilon}) \cdot \nabla_v f_\epsilon = 0 \\ f_{t=0} = f_0 \end{cases}$$

Frénod and Sonnendrücker proved the following theorem:

Theorem 1.1. *For each ϵ , let f_ϵ be the unique solution of the scaled Vlasov equation in $L_t^\infty(L_{x,v}^1 \cap L_{x,v}^2)$. Then the following convergence holds as ϵ tends to 0:*

$$f_\epsilon \rightharpoonup f \text{ weak-}^* L_t^\infty(L_{x,v}^2) \quad (1.13)$$

where $f \in L_t^\infty(L_{x,v}^2)$ is the unique solution to:

$$\begin{aligned} \partial_t f + v_\parallel \cdot \nabla_x f + \frac{1}{2\pi} \left(\int_0^{2\pi} \mathcal{R}(\tau) E(t, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_x f \\ + \frac{1}{2\pi} \left(\int_0^{2\pi} R(\tau) E(t, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_v f = 0 \\ f|_{t=0} = \frac{1}{2\pi} \left(\int_0^{2\pi} f_0(x + \mathcal{R}(\tau)v, R(\tau)v) d\tau \right) \end{aligned}$$

denoting by R and \mathcal{R} the linear operators defined by:

$$R(\tau) = \begin{bmatrix} \cos \tau & -\sin \tau & 0 \\ \sin \tau & \cos \tau & 0 \\ 0 & 0 & 1 \end{bmatrix}, \mathcal{R}(\tau) = (-R(-\pi/2) + R(-\pi/2 + \tau))$$

2. The pseudo 2D case:

The Vlasov-Poisson system considered in this case is the following 2D system:

$$\partial_t f_\epsilon + \frac{v}{\epsilon} \cdot \nabla_x f_\epsilon + \left(E_\epsilon + \frac{v^\perp}{\epsilon} \right) \cdot \nabla_v f_\epsilon = 0 \quad (1.14)$$

$$f_\epsilon|_{t=0} = f_0 \quad (1.15)$$

$$E_\epsilon = -\nabla V_\epsilon, -\Delta_x V_\epsilon = \rho_\epsilon \quad (1.16)$$

$$\rho_\epsilon = \int f_\epsilon dv \quad (1.17)$$

If $v = (v_x, v_y)$, v^\perp is defined by $(v_y, -v_x)$.

We recall that there exist global weak solutions of Vlasov-Poisson systems in the sense of Arsenev ([3]).

Assuming here that $f_0 \geq 0$, $f_0 \in L^1_{x,v} \cap L^p_{x,v}$ (for some $p > 2$) and that the initial energy is bounded, Frénod and Sonnendrücker proved the following theorem (we voluntarily write an unprecise meta-version of the result)

Theorem 1.2. *For each ϵ , let (f_ϵ, E_ϵ) be a solution in the sense of Arsenev to (1.14)-(1.17).*

Then, up to a subsequence, f_ϵ weakly converges to a function f . Moreover, there exists a function G such that :

$$f = \int_0^{2\pi} G(t, x + \mathcal{R}(\tau)v, R(\tau)v) d\tau \quad (1.18)$$

and G satisfies :

$$\partial_t G + \frac{1}{2\pi} \left(\int_0^{2\pi} \mathcal{R}(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_x G$$

$$+ \frac{1}{2\pi} \left(\int_0^{2\pi} R(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_v G = 0$$

$$G|_{t=0} = f_0$$

$$\mathcal{E} = -\nabla \Phi, \quad -\Delta \Phi = \int G(t, x + \mathcal{R}(\tau)v, R(\tau)v) dv$$

denoting by R and \mathcal{R} the linear operators defined by :

$$R(\tau) = \begin{bmatrix} \cos \tau & -\sin \tau \\ \sin \tau & \cos \tau \end{bmatrix}, \mathcal{R}(\tau) = (R(-\pi/2) - R(-\pi/2 + \tau))$$

In this case, we have to introduce an additional variable, the ‘‘fast-time’’ variable τ which comes from the fact that we need to precisely describe the oscillations in order to study the limit in non-linear terms.

Note that the authors actually developed a generic framework that allows them to deal with different scalings and to give a precise approximation at any order. We do not wish to do so in our study.

1.3 A bit of homogenization theory and some useful definitions

Let us now precisely state the “2-scale” convergence tools used in this paper.

Definition. Let X be a separable Banach space, X' be its topological dual space and (\cdot, \cdot) the duality bracket between X' and X . For all $\alpha > 0$, denote by $\mathcal{C}_\alpha(\mathbb{R}, X)$ (respectively $L_\alpha^{q'}(\mathbb{R}; X')$) the space of α -periodic continuous (respectively $L^{q'}$) functions on \mathbb{R} with values in X . Let $q \in [1; \infty[$.

Given a sequence (u_ϵ) of functions belonging to the space $L^{q'}(0, t; X')$ and a function $U^0(t, \theta) \in L^{q'}(0, T; L_\alpha^{q'}(\mathbb{R}; X'))$ we say that

$$u_\epsilon \text{ 2-scale converges to } U^0$$

if for any function $\Psi \in L^q(0, T; \mathcal{C}_\alpha(\mathbb{R}, X))$ we have:

$$\lim_{\epsilon \rightarrow 0} \int_0^T \left(u_\epsilon(t), \Psi \left(t, \frac{t}{\epsilon} \right) \right) dt = \frac{1}{\alpha} \int_0^T \int_0^\alpha (U^0(t, \tau), \Psi(t, \tau)) d\tau dt \quad (1.19)$$

Theorem 1.3. Given a sequence (u_ϵ) bounded in $L^{q'}(0, t; X')$, there exists for all $\alpha > 0$ a function $U_\alpha^0 \in L^{q'}(0, T; L_\alpha^{q'}(\mathbb{R}; X'))$ such that up to a subsequence,

$$u_\epsilon \text{ 2-scale converges to } U_\alpha^0$$

The profile U_α^0 is called the α -periodic two scale limit of u_ϵ and the link between U_α^0 and the weak- $*$ limit u of u_ϵ is given by:

$$\frac{1}{\alpha} \int_0^\alpha U^0 d\tau = u \quad (1.20)$$

We also introduce some notations:

Notations. We define for all $p \in [1; \infty]$ the space $L_{x,v}^p := L_x^p(\mathbb{T}^d, (L_v^p(\mathbb{R}^d)))$.

In the same fashion, we define the spaces $L_{t,x}^p, L_{t,x,v}^p \dots$

Let $L_{2\pi,\tau}^p$ be the space of 2π -periodic functions of τ which are in L_τ^p .

Let $L_{x,loc}^p$ be the space of functions f such that for all infinitely differentiable cut-off functions $\varphi \in \mathcal{C}_c^\infty$, φf belongs to L_x^p . We will say that a sequence (f_ϵ) is uniformly bounded in $L_{x,loc}^p$ if for each compact set K , the sequence of the restrictions to K is uniformly bounded in L_x^p with respect to ϵ (but this bound can depend on K).

We will also use the same notations for Sobolev spaces $W^{s,p}$ ($s \in \mathbb{R}$).

1.4 Statement of the result

In this paper we prove that the 2-scale convergence established in the previous 2D case is also true in our 3D framework. The difficulty comes from the fact that there is no uniform elliptic regularity for the electric field because of the factor ϵ^2 in front of Δ_{x_\parallel} in the Poisson equation:

$$V_\epsilon - \epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon = \int f_\epsilon dv - \int f_\epsilon dv dx$$

In particular there is no a priori regularity on x_\parallel and therefore no strong compactness. Nevertheless, we actually prove that due to the particular form of the asymptotic equation, the moments of the solution with respect to v_\parallel are more regular in x_\parallel than the solution itself. We can then easily pass to the weak limit.

The reason why we have opted for this strange Poisson equation instead of the usual one will appear at the end of the next section and especially in the last one. Roughly speaking it allows us to “kill” plasma waves which appear in the parallel direction due to the quasineutrality.

Notice that this result is in the same spirit as the proof of the weak stability of the Vlasov-Maxwell system by DiPerna and Lions ([8]), where the authors have regularity on moments, by opposition to the proof of the weak stability of the Vlasov-Poisson system by Arsenev ([3]), where the author has compactness on the electric field. Actually our result is a kind of a hybrid one, since we get on one hand regularity with respect to x_\perp by elliptic regularity and in the other hand regularity with respect to x_\parallel by averaging.

We assume here that the initial data $(f_{\epsilon,0})_{\epsilon>0}$ satisfy the following conditions:

- $f_{\epsilon,0} \geq 0$ (positivity)
- $(f_{\epsilon,0})_{\epsilon>0}$ is uniformly bounded with respect to ϵ in $L^1_{x,v} \cap L^p_{x,v}$ (for some $p > 3$) and for each ϵ , $\int f_{\epsilon,0} dx dv = 1$.
- The initial energy is uniformly bounded with respect to ϵ :

$$\left(\int f_{\epsilon,0} |v|^2 dv dx + \epsilon \int V_{\epsilon,0}^2 dx + \epsilon \int |\nabla_{x_\perp} V_{\epsilon,0}|^2 dx + \epsilon^3 \int |\nabla_{x_\parallel} V_{\epsilon,0}|^2 dx \right) \leq C$$

Theorem 1.4. *For each ϵ , let (f_ϵ, E_ϵ) in $L_t^\infty(L^1_{x,v} \cap L^p_{x,v}) \times L_t^\infty(L^2_x)$ be a global weak solution in the sense of Arsenev to (1.12). Then up to a subsequence we have the following convergence as ϵ tends to 0:*

$$f_{\epsilon,0} \text{ weakly-}^* \text{ converges to } f_0 \in L^p_{x,v} \quad (1.21)$$

$$f_\epsilon \text{ 2-scale converges to } F \in L_t^\infty(L^{2\pi,\tau}(L^1_{x,v} \cap L^p_{x,v})) \quad (1.22)$$

$$E_\epsilon \text{ 2-scale converges to } \mathcal{E} \in L_t^\infty(L^{2\pi,\tau}(L^{3/2}_{x_\parallel}(W^{1,\frac{3}{2}}_{x_\perp}))) \quad (1.23)$$

Moreover, there exists a function $G \in L_t^\infty(L^1_{x,v} \cap L^p_{x,v})$ such that:

$$F(t, \tau, x, v) = G(t, x + \mathcal{R}(\tau)v, R(\tau)v) \quad (1.24)$$

and (G, \mathcal{E}) is solution to:

$$\begin{aligned} \partial_t G + v_\parallel \cdot \nabla_x G + \frac{1}{2\pi} \left(\int_0^{2\pi} \mathcal{R}(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_x G \\ + \frac{1}{2\pi} \left(\int_0^{2\pi} R(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_v G = 0 \end{aligned}$$

$$G|_{t=0} = f_0$$

$$\mathcal{E} = (-\nabla_\perp V, 0), \quad V - \Delta_\perp V = \int G(t, x + \mathcal{R}(\tau)v, R(\tau)v) dv - \int f_0 dv dx$$

denoting by R and \mathcal{R} the linear operators defined by:

$$R(\tau) = \begin{bmatrix} \cos \tau & -\sin \tau & 0 \\ \sin \tau & \cos \tau & 0 \\ 0 & 0 & 1 \end{bmatrix}, \quad \mathcal{R}(\tau) = (R(-\pi/2) - R(-\pi/2 + \tau))$$

As it has been said, for the proof of this theorem, we will first prove a proposition which gives the regularity of moments in v_{\parallel} of the solution. For this, we use an averaging lemma. The beginning of the proof is very similar to the proof in the $2D$ case, but we will give it again for the sake of completeness.

Remarks. 1. The assumption on the initial energy may, at first sight, look a bit restrictive but in the “usual” Vlasov-Poisson scaling, it only means that the initial electric potential and field are bounded in L^2 .

2. The constant $q = 3$ will come quite naturally from Lemma 2.3 and Proposition 3.1.

3. This theorem implies that for a given non-negative initial data $G|_{t=0} = G_0$ in $L^1_{x,v} \cap L^p_{x,v}$ (with $p > 3$) and satisfying the energy bound, the asymptotic system admits at least one global weak solution $G \in L^\infty_t(L^1_{x,v} \cap L^p_{x,v})$. With the additional assumptions on the initial data:

$$\begin{aligned} G_0 &\in W^{1,1}_{x,v}, \\ \|(1 + |v|^4)G_0\|_{L^\infty_{x,v}} &< \infty, \\ \|(1 + |v|^4)DG_0\|_{L^\infty_{x,v}} &< \infty. \end{aligned}$$

we are actually able to prove the uniqueness of the solution, using the same ideas than Degond in [7] (and also used afterwards by Saint-Raymond in a gyrokinetic context ([21])). Hence, it means that if the whole sequence $(f_{\epsilon,0})$ satisfies the same additional estimates, uniformly with respect to ϵ , and weak $*$ converges to some f_0 then there is also convergence for the whole sequence (f_ϵ) .

2 A priori uniform estimates for the scaled Vlasov-Poisson system

2.1 Conservation of L^p norms and energy for the scaled system

In this section we give a priori estimates which are very classical for the Vlasov-Poisson system (used for example in [11], [12], [15]). In order to recall how one can get them, we will give some formal computations. If one wants to have rigorous proofs, one should deal with smooth and compactly supported functions, namely with a sequence $(f_\epsilon^n)_{n \geq 0}$ of solutions of some regularized Vlasov-Poisson equations then pass to the limit (that is the way one can classically build a global weak solution in the sense of Arsenev ([3])).

First, as usual for such Vlasov equations, L^p norms are conserved (we work here at a fixed ϵ):

Lemma 2.1. *For all $1 \leq p \leq \infty$,*

$$\forall t \geq 0, \|f(t)\|_{L^p_{x,v}} \leq \|f(0)\|_{L^p_{x,v}} \quad (2.1)$$

Moreover, $f_0 \geq 0$ if and only if $\forall t \geq 0, f(t) \geq 0$ (referred to as the maximum principle)

That precisely means that if $f_0 \in L^p_{x,v}$, then $f \in L^\infty_t(L^p_{x,v})$.

Let us now compute the energy for the scaled system:

Lemma 2.2. *We have the estimate:*

$$\mathcal{E}_\epsilon(t) = \left(\int f_\epsilon |v|^2 dv dx + \epsilon \int V_\epsilon^2 dx + \epsilon \int |\nabla_{x_\perp} V_\epsilon|^2 dx + \epsilon^3 \int |\nabla_{x_\parallel} V_\epsilon|^2 dx \right) \leq \mathcal{E}_\epsilon(0) \quad (2.2)$$

In particular if there exists $C > 0$ independent of ϵ such that $\mathcal{E}_\epsilon(0) \leq C$, then:

$$\int f_\epsilon |v|^2 dv dx \leq C \quad (2.3)$$

Formal proof. We multiply the scaled Vlasov equation by $|v|^2$ and integrate with respect to x and v .

$$\int \partial_t f_\epsilon |v|^2 dv dx + \int E_\epsilon \cdot \nabla_v f_\epsilon |v|^2 dv dx = \frac{d}{dt} \left(\int f_\epsilon |v|^2 dv dx \right) - 2 \int E_\epsilon(x) \cdot v f_\epsilon dv dx = 0$$

We then integrate the Vlasov equation with respect to v . We get the so called conservation of charge:

$$\frac{d}{dt} \left(\int f dv \right) + \nabla_{x_\parallel} \cdot \left(\int f v_\parallel dv \right) + \frac{\nabla_{x_\perp}}{\epsilon} \cdot \left(\int f v_\perp dv \right) = 0 \quad (2.4)$$

Therefore, we have:

$$\begin{aligned} \int E_\epsilon(x) \cdot v f_\epsilon dv dx &= - \int (\nabla_{x_\perp} V_\epsilon, \epsilon \nabla_{x_\parallel} V_\epsilon) \cdot v f_\epsilon dv dx \\ &= \int V_\epsilon \left(\nabla_{x_\perp} \cdot (f_\epsilon v_\perp) + \epsilon \nabla_{x_\parallel} \cdot (f_\epsilon v_\parallel) \right) dv dx \\ &= -\epsilon \int V_\epsilon \partial_t f_\epsilon dv dx \end{aligned}$$

Finally, using the Poisson equation, we get:

$$\begin{aligned} -\epsilon \int V_\epsilon \partial_t f_\epsilon dv dx &= -\epsilon \int V_\epsilon \partial_t \left(V_\epsilon - \epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon \right) dx \\ &= -\epsilon \left(\int V_\epsilon \partial_t V_\epsilon dx + \int \nabla_{x_\perp} V_\epsilon \partial_t \nabla_{x_\perp} V_\epsilon dx + \epsilon^2 \int \nabla_{x_\parallel} \partial_t V_\epsilon \nabla_{x_\parallel} V_\epsilon dx \right) \\ &= -\epsilon \frac{1}{2} \frac{d}{dt} \left(\int V_\epsilon^2 dx + \int |\nabla_{x_\perp} V_\epsilon|^2 dx + \epsilon^2 \int |\nabla_{x_\parallel} V_\epsilon|^2 dx \right) \end{aligned}$$

Thus it comes:

$$\frac{d}{dt} \left(\int f_\epsilon |v|^2 dv dx + \epsilon \int V_\epsilon^2 dx + \epsilon \int |\nabla_{x_\perp} V_\epsilon|^2 dx + \epsilon^3 \int |\nabla_{x_\parallel} V_\epsilon|^2 dx \right) = 0 \quad (2.5)$$

□

2.2 Regularity of the electric field

Let us recall a classical lemma obtained by a standard real interpolation argument:

Lemma 2.3. *Let $f(x, v)$ be a measurable positive function on $\mathbb{R}^3 \times \mathbb{R}^3$. Then:*

$$\int \left(\int f(x, v) dv \right)^{3/2} dx \leq C \|f\|_{L_{x,v}^3}^{3/4} \left(\int |v|^2 f dx dv \right)^{3/4} \quad (2.6)$$

Proof. For any $R > 0$, we can write the following decomposition:

$$\begin{aligned}\int f_\epsilon dv &= \int_{|v| \leq R} f dv + \int_{|v| > R} f dv \\ &\leq CR^2 \|f\|_{L_v^3} + \frac{1}{R^2} \int |v|^2 f dv\end{aligned}$$

Then we can take R such that $R^2 \|f\|_{L_v^3} = \frac{1}{R^2} \int |v|^2 f dv$ so that we get:

$$\int f dv \leq C \left(\int f^3 dv \right)^{1/6} \left(\int |v|^2 f dv \right)^{1/2} \quad (2.7)$$

We then raise the quantities to the power $3/2$, integrate with respect to x and use Hölder's inequality which gives the estimate. \square

By conservation of the L^3 norm and the uniform bound on the initial energy, Lemmas 2.2 and 2.3 entail that:

$$\rho_\epsilon \in L_t^\infty(L_x^{3/2}) \quad (2.8)$$

and the norm is bounded uniformly with respect to ϵ .

We now use the Poisson equation to compute the regularity of the electric field. Let us recall that:

$$\begin{aligned}E_\epsilon &= \left(-\epsilon \nabla_{x_\parallel} V_\epsilon, -\nabla_{x_\perp} V_\epsilon \right) \\ V_\epsilon - \epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon &= \rho_\epsilon - \int \rho_0 dx\end{aligned}$$

Lemma 2.4. *With the above notations and assumptions:*

E_ϵ is uniformly bounded with respect to ϵ in $L_t^\infty(L_{x_\parallel}^{3/2}(W_{x_\perp}^{1,3/2}))$

Proof. Let $\epsilon > 0$ and $t > 0$ be fixed. For the sake of simplicity we write V instead of V_ϵ and E instead of E_ϵ .

For any function $f(x_\parallel, x_\perp)$, define the rescaled function $\tilde{f}(z, x_\perp)$ by

$$\tilde{f}\left(\frac{x_\parallel}{\epsilon}, x_\perp\right) = \epsilon^{\frac{2}{3}} f(x_\parallel, x_\perp)$$

so that:

$$\|\tilde{f}(z, x_\perp)\|_{L_z^{3/2}} = \|f(x_\parallel, x_\perp)\|_{L_{x_\parallel}^{3/2}} \quad (2.9)$$

The Poisson equation becomes:

$$\tilde{V} - \Delta_z \tilde{V} - \Delta_{x_\perp} \tilde{V} = \tilde{\rho} - \epsilon^{\frac{2}{3}} \int \rho_0 dx$$

and the scaled electric field is given by:

$$\tilde{E} = \left(-\nabla_z \tilde{V}, -\nabla_{x_\perp} \tilde{V} \right)$$

Since $\rho_\epsilon(t, \cdot, \cdot)$ and V_ϵ are uniformly bounded in $L_x^{3/2}$, standard results of elliptic regularity on the torus $\mathbb{T}^2 \times \frac{1}{\epsilon} \mathbb{T}$ show that there exists $C > 0$ independent of ϵ such that:

$$\|\tilde{V}\|_{W_{z, x_\perp}^{2, 3/2}} \leq C \left\| \tilde{\rho} - \epsilon^{\frac{2}{3}} \int \rho_0 dx \right\|_{L_{z, x_\perp}^{3/2}}$$

Remark. Notice here that due to the dilatation of order $\frac{1}{\epsilon}$ in the parallel direction, being periodic in this direction does not make things easier.

Thanks to (2.9) we get:

$$\|\tilde{V}\|_{W_{z,x_\perp}^{2,3/2}} \leq C \left\| \rho - \int \rho_0 dx \right\|_{L_t^\infty(L_{x_\perp}^{3/2})} \leq C_0$$

with C_0 independent of ϵ .

Consequently, we have:

$$\|\tilde{E}\|_{L_z^{3/2}(W_{x_\perp}^{1,3/2})} \leq \|\tilde{E}\|_{W_z^{1,3/2}(W_{x_\perp}^{1,3/2})} \leq C_0$$

Finally from (2.9) we get

$$\|E_\epsilon\|_{L_t^\infty(L_{x_\parallel}^{3/2}(W_{x_\perp}^{1,3/2}))} \leq C_0$$

□

We can see as expected that the regularity of the electric field with respect to the x_\parallel variable is not sufficient to get some strong compactness.

Remarks. 1. We can write the identity:

$$-\Delta_{x_\perp} V_\epsilon = -\Delta_{x_\perp} (Id - \epsilon^2 \Delta_{x_\parallel} - \Delta_{x_\perp})^{-1} \left(\rho_\epsilon - \int \rho_\epsilon dx \right) \quad (2.10)$$

so that, thanks to elliptic estimates on the torus \mathbb{T}^2 , $V_\epsilon \in L_{x_\parallel}^{3/2}(W_{x_\perp}^{2,3/2})$. Consequently, $\partial_{x_\parallel} V_\epsilon$ is bounded in $W_{x_\parallel}^{-1,3/2}(W_{x_\perp}^{2,3/2})$. In other words, $E_{\epsilon,\parallel} = -\epsilon \partial_{x_\parallel} V_\epsilon$ tends to zero in the sense of distributions.

2. A typical function φ_ϵ such that φ_ϵ is bounded in L^p and $\frac{1}{\epsilon} \varphi_\epsilon$ is bounded in $W^{-1,p}$ is the oscillating function $\cos(\frac{1}{\epsilon}x)$. This indicates that $E_{\epsilon,\parallel}$ oscillates with a frequency of order $\frac{1}{\epsilon}$ in the parallel direction.
3. If we work with the usual Poisson equation

$$-\epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon = \rho_\epsilon - \int \rho_\epsilon dx$$

we only get homogeneous estimates for V_ϵ and we have not been able to deal with such anisotropic estimates in the following of the paper (namely in the estimates of Proposition 3.1). Roughly speaking, if V is a solution of the Poisson equation $-\Delta V = \rho$ with $\rho \in L^{3/2}(\mathbb{R}^3)$, we can only say that $V \in \dot{W}^{2,3/2}(\mathbb{R}^3)$ (the homogeneous Sobolev space) and not $W^{2,3/2}$.

4. This difficulty seems to be not only a technical one, but appears to be linked to the existence of plasma waves (with frequency and magnitude of order $\frac{1}{\sqrt{\epsilon}}$) in the parallel direction which prevents us from passing directly to the limit $\epsilon \rightarrow 0$ (see [19] and last section).

3 Proof of Theorem 1.4

Proof. The first two steps are identical to the one given in [12]. For the sake of completeness we recall here the main arguments and refer to [12] for the details.

Step 1: Deriving the constraint equation

First of all, since (f_ϵ) is bounded in $L_t^\infty(L_{x,v}^1 \cap L_{x,v}^p)$, Theorem 1.3 shows that for all $\alpha > 0$:

$$f_\epsilon \text{ 2-scale converges to } F_\alpha \in L^\infty(0, T; L_\alpha^\infty(\mathbb{R}; L_{x,v}^p))$$

Let $\Psi(t, \tau, x, v)$ be an α -periodic oscillating test function in τ and define:

$$\Psi^\epsilon \equiv \Psi\left(t, \frac{t}{\epsilon}, x, v\right)$$

We start by writing the weak formulation of the scaled Vlasov equation against Ψ^ϵ . Since

$$\nabla_{x_\parallel} \cdot v_\parallel = \nabla_{x_\perp} \cdot v_\perp = \operatorname{div}_v \left(E_\epsilon + \frac{v \wedge e_z}{\epsilon} \right) = 0,$$

we get the following equation:

$$\begin{aligned} \int f_\epsilon \left((\partial_t \Psi)^\epsilon + \frac{1}{\epsilon} (\partial_\tau \Psi)^\epsilon + v_\parallel \cdot (\nabla_x \Psi)^\epsilon + \frac{v_\perp}{\epsilon} \cdot (\nabla_x \Psi)^\epsilon + \left(E_\epsilon + \frac{v \wedge e_z}{\epsilon} \right) \cdot (\nabla_v \Psi)^\epsilon \right) dt dx dv \\ = - \int f_0 \Psi(0, 0, x, v) dx dv \end{aligned}$$

Multiply then by ϵ and pass up to a subsequence to the (2-scale) limit. We get the so called constraint equation for the α -periodic profile F_α :

$$\partial_\tau F_\alpha + v_\perp \cdot \nabla_x F_\alpha + v \wedge e_z \cdot \nabla_v F_\alpha = 0, \quad (3.1)$$

which means that F_α is constant along the characteristics:

$$\frac{dV}{d\tau} = V \wedge e_z \quad (3.2)$$

$$\frac{dX}{d\tau} = V_\perp \quad (3.3)$$

A straightforward calculation therefore shows that there exists $F_\alpha^0 \in L^\infty(0, T, L_{x,v}^p)$ such that:

$$F_\alpha(t, \tau, x, v) = F_\alpha^0(t, x + \mathcal{R}(\tau)v, R(\tau)v) \quad (3.4)$$

with:

$$R(\tau) = \begin{bmatrix} \cos \tau & -\sin \tau & 0 \\ \sin \tau & \cos \tau & 0 \\ 0 & 0 & 1 \end{bmatrix}, \mathcal{R}(\tau) = (-R(-\pi/2) + R(-\pi/2 + \tau))$$

$$\text{i.e. } \mathcal{R}(\tau) = \begin{bmatrix} \sin \tau & \cos \tau - 1 & 0 \\ 1 - \cos \tau & \sin \tau & 0 \\ 0 & 0 & 0 \end{bmatrix}.$$

Since R and \mathcal{R} are 2π -periodic, we will consider the 2π profile: indeed if α and 2π were incommensurable, F_α could not depend on τ and consequently we would have no information on the oscillations.

Step 2: Filtering the essential oscillation

We now look for the equation satisfied by $F_{2\pi}^0 := G$; we introduce the filtered function g_ϵ :

$$g_\epsilon(t, x, v) = f_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v, R(-t/\epsilon)v) \quad (3.5)$$

(meaning that we have removed the oscillations)

We easily compute the equation satisfied by g_ϵ :

$$\begin{aligned} \partial_t g_\epsilon + v_{\parallel} \cdot \nabla_x g_\epsilon + \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x g_\epsilon \\ + R(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v g_\epsilon = 0 \end{aligned} \quad (3.6)$$

Remark. Note here that g_ϵ 2-scale converges to G , and since it does not depend on τ , it also weakly converges to G .

Step 3: Getting some regularity on moments

From now on, the goal is to get some compactness for the moments of g_ϵ with respect to v_{\parallel} . The main tool is the following averaging lemma proved by Bézard in [6], which is a refined version of the fundamental result of DiPerna, Lions and Meyer ([9]):

Theorem 3.1. *Let $1 < p \leq 2$. Let $f, g \in L^p(dt \otimes dx \otimes dv)$ be solutions of the following transport equation*

$$\partial_t f + v \cdot \nabla_x f = (I - \Delta_{t,x})^{\tau/2} (I - \Delta_v)^{m/2} g \quad (3.7)$$

with $m \in \mathbb{R}^+$, $\tau \in [0, 1[$. Then $\forall \Psi \in \mathcal{C}_c^\infty(\mathbb{R}^d)$, $\rho_\Psi(t, x) = \int f(t, x, v) \Psi(v) dv \in W_{t,x}^{s,p}(\mathbb{R} \times \mathbb{R}^d)$ where

$$s = \frac{1 - \tau}{(1 + m)p'} \quad (3.8)$$

Moreover,

$$\|\rho_\Psi\|_{W_{t,x}^{s,p}(\mathbb{R} \times \mathbb{R}^d)} \leq C \left(\|f\|_{L^p(dt \otimes dx \otimes dv)} + \|(I - \Delta_{t,x})^{\tau/2} (I - \Delta_v)^{m/2} g\|_{W_{t,x}^{-\tau,p}(W_v^{-m,p})} \right) \quad (3.9)$$

(C is a positive constant independent of f and g)

Averaging lemmas are an important feature of transport equations: since the transport equation (3.7) is hyperbolic, one can obviously not expect the solution f to be more regular than the right hand side or the initial data. Nevertheless, if one considers the averaged quantity ρ_Ψ , one can actually notice a gain of regularity. This phenomenon was first observed independently by Golse, Perthame and Sentis ([14]) and Agoshkov ([1]) then was formulated in a precise way for the first time by Golse, Lions, Perthame and Sentis (see [13]); it is referred to as “velocity averaging”. There exists many refined versions of these results and numerous interesting applications in kinetic theory, but we shall not dwell on that. We simply point out that this tool has been successfully applied to Vlasov equations, for instance to prove the existence of global weak solutions to the Vlasov-Maxwell system as it has been done by DiPerna and Lions ([8]).

Remarks. 1. The fundamental averaging lemmas proved by DiPerna, Lions and Meyer ([9]), which lead to a Besov estimate instead of a Sobolev estimate, would have been sufficient for our purpose, since we do not need the optimal regularity.

2. These lemmas have been proved for functions with values in \mathbb{R} but the proof is the same with functions with values in more complicated Banach spaces. Indeed one can consider “abstract” Banach norms instead of the usual norm in \mathbb{R} and the proof is then formally identical.

Equipped with this tool, we can now prove that moments in v_{\parallel} are more regular with respect to t and x_{\parallel} than the solution itself.

Proposition 3.1. *For each $\epsilon > 0$, let g_{ϵ} be a function in $L^1_{x,v} \cap L^p_{x,v}$ (with $p > 3$) bounded uniformly with respect to ϵ and satisfying:*

$$\begin{aligned} \partial_t g_{\epsilon} + v_{\parallel} \cdot \nabla_x g_{\epsilon} + \mathcal{R}(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x g_{\epsilon} \\ + \mathcal{R}(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v g_{\epsilon} = 0 \end{aligned}$$

with E_{ϵ} uniformly bounded in $L^{\infty}_t(L^{3/2}_{x_{\parallel}}(W^{1,3/2}_{x_{\perp}}))$.

Let $\Psi \in \mathcal{D}(\mathbb{R})$. Define

$$\eta_{\epsilon}(t, x, v_{\perp}) = \int g_{\epsilon}(t, x, v) \Psi(v_{\parallel}) dv_{\parallel}$$

Then,

$$\eta_{\epsilon} \text{ is uniformly bounded in } W^{s,\gamma}_{t,loc}(W^{s,\gamma}_{x_{\parallel},loc}(W^{-1,\gamma}_{x_{\perp},v_{\perp},loc})) \quad (3.10)$$

for $\gamma \in]1; 2[$ defined by $\frac{1}{\gamma} = \frac{2}{3} + \frac{1}{p}$ and some $s \in]0; 1[$ (depending on γ)

Proof. • The first step is to localize the equation. Let K be the cartesian product of compact sets:

$$K = [0, T] \times K_{x_{\parallel}} \times K_{x_{\perp}} \times K_{v_{\parallel}} \times K_{v_{\perp}}$$

We now consider some positive smooth function $\Phi(t, x_{\parallel}, x_{\perp}, v_{\parallel}, v_{\perp})$ which is \mathcal{C}_c^{∞} and which satisfies the condition:

$$\Phi \equiv 0 \text{ outside } K \quad (3.11)$$

The equation satisfied by $g_{\epsilon}\Phi$ is the following one:

$$\begin{aligned} \partial_t(g_{\epsilon}\Phi) + v_{\parallel} \cdot \nabla_x(g_{\epsilon}\Phi) = & \underbrace{-\mathcal{R}(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x(g_{\epsilon}\Phi)}_{(1)} \\ & - \underbrace{R(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v(g_{\epsilon}\Phi)}_{(2)} - \partial_t(\Phi)g_{\epsilon} - v_{\parallel} \cdot \nabla_x(\Phi)g_{\epsilon} \\ & + \mathcal{R}(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x(\Phi)g_{\epsilon} + R(t/\epsilon) E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v(\Phi)g_{\epsilon} \end{aligned}$$

The idea is now to consider this equation as a kinetic equation with respect to the variables $(t, x_{\parallel}, v_{\parallel})$ and with values in an abstract Banach space (which will be $W^{-1,\gamma}_{x_{\perp},v_{\perp}}$). We then only study the first two terms of the right-hand side (noticing that the other terms have more regularity than these ones).

From now on, for the sake of simplicity and readability, we will write L^p and $W^{s,p}$ norms without always specifying that they are taken on the compact support of Φ .

• **Estimate on the first term (1)**

Since E_{ϵ} does not depend on v , we have:

$$E_{\epsilon} \in L^{\infty}_t(L^{3/2}_{x_{\parallel}}(W^{1,\infty}_v(W^{1,3/2}_{x_{\perp}})))$$

In particular if we restrict to compact supports:

$$E_\epsilon \in L_t^{3/2} \left(L_{x_\parallel}^{3/2} \left(W_{v_\parallel}^{1,3/2} \left(W_{v_\perp}^{1,3/2} \left(W_{x_\perp}^{1,3/2} \right) \right) \right) \right)$$

The second point is that the differential operator

$$\mathcal{R}(t/\epsilon)E_\epsilon(t, x + \mathcal{R}(t/\epsilon)v) \cdot \nabla_x$$

involves only derivatives with respect to the x_\perp variable and not in the parallel direction (this remark is fundamental for using an averaging lemma). Product laws in Sobolev spaces give then:

$$\begin{aligned} & \|\mathcal{R}(t/\epsilon)E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x (g_\epsilon \Phi)\|_{W_{x_\perp, v_\perp}^{-1, \gamma} (K_{x_\perp} \times K_{v_\perp})} \leq \\ & \|E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v)\|_{W_{x_\perp, v_\perp}^{1, 3/2} (K_{x_\perp} \times K_{v_\perp})} \|\mathcal{R}(t/\epsilon) \cdot \nabla_x (g_\epsilon \Phi)\|_{W_{x_\perp, v_\perp}^{-1, p} (K_{x_\perp} \times K_{v_\perp})} \end{aligned}$$

where $\frac{1}{\gamma} = \frac{2}{3} + \frac{1}{p}$.

Remark. The regularity of (1) with respect to v_\perp is not optimal (since it involves no derivative in v_\perp for g_ϵ). Nevertheless we are interested in the regularity of the whole right hand side, and we will see that the term (2) has this regularity in v_\perp .

The change of variables $(x, v) \mapsto (x + \mathcal{R}(s)v, \mathcal{R}(s)v)$ has unit Jacobian for all $s \in \mathbb{R}$ so that:

$$\|E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v)\|_{W_{x_\perp, v_\perp}^{1, 3/2} (K_{x_\perp} \times K_{v_\perp})} = \|E_\epsilon(t, x)\|_{W_{x_\perp, v_\perp}^{1, 3/2} (J_{x_\perp} \times J_{v_\perp})}$$

where J_{x_\perp} and J_{v_\perp} are the compact images of K_{x_\perp} and K_{v_\perp} by the change of variables. Moreover, one gets:

$$\|\mathcal{R}(t/\epsilon) \cdot \nabla_x (g_\epsilon \Phi)\|_{W_{x_\perp, v_\perp}^{-1, p}} \leq C \|g_\epsilon \Phi\|_{L_{x_\perp, v_\perp}^p}$$

Finally we have after integrating in $t, x_\parallel, v_\parallel$ and thanks to Hölder's inequality:

$$\begin{aligned} & \|\mathcal{R}(t/\epsilon)E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x (g_\epsilon \Phi)\|_{L_{t, x_\parallel, v_\parallel}^\gamma (W_{x_\perp, v_\perp}^{-1, \gamma})} \\ & \leq C \|E_\epsilon(t, x)\|_{L_{t, x_\parallel, v_\parallel}^{3/2} (W_{x_\perp, v_\perp}^{1, 3/2})} \|g_\epsilon \Phi\|_{L_{t, x_\parallel, v_\parallel}^p (L_{x_\perp, v_\perp}^p)} \end{aligned}$$

and C is a constant independant of ϵ .

- **Estimate on the second term (2)**

By the same method one gets:

$$\begin{aligned} & \|R(t/\epsilon)E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v (g_\epsilon \Phi)\|_{L_{t, x_\parallel}^\gamma (W_{v_\parallel}^{-1, \gamma} (W_{x_\perp, v_\perp}^{-1, \gamma}))} \\ & \leq C \|E_\epsilon(t, x)\|_{L_{t, x_\parallel}^{3/2} (W_{x_\perp, v}^{1, 3/2})} \|g_\epsilon \Phi\|_{L_{t, x_\parallel}^p (L_{x_\perp, v}^p)} \end{aligned}$$

Finally we see that the right hand side is uniformly bounded in:

$$L_{t, x_\parallel, \text{loc}}^\gamma (W_{v_\parallel, \text{loc}}^{-1, \gamma} (W_{x_\perp, v_\perp, \text{loc}}^{-1, \gamma}))$$

- **Regularity of the moments**

By Theorem 3.1 (and the second remark given afterwards), for all $\Psi \in \mathcal{C}_c^\infty$, the moment:

$$\eta_\epsilon(t, x, v_\perp) = \int g_\epsilon(t, x, v) \Psi(v_\parallel) dv_\parallel$$

is then uniformly bounded in the space $W_{t,x_\parallel, \text{loc}}^{s,\gamma}(W_{x_\perp, v_\perp, \text{loc}}^{-1,\gamma})$ with $s > 0$ given by $s = \frac{1}{2\gamma}$. □

We can now prove that the sequence of moments η_ϵ is compact in a space of distributions which is the dual of some space where the sequence (E_ϵ) is uniformly bounded.

Corollary 3.1. *There exists $\theta \in]0, 1[$ and $\eta \in W_{t, \text{loc}}^{s\theta, 3}(W_{x_\parallel, \text{loc}}^{s\theta, 3}(W_{x_\perp, v_\perp, \text{loc}}^{-\theta, 3}))$ such that for all $\xi > 0$, up to a subsequence:*

$$\eta_\epsilon \rightarrow \eta \text{ strongly in } L_{t, \text{loc}}^3(L_{x_\parallel, \text{loc}}^3(W_{x_\perp, v_\perp, \text{loc}}^{-\theta-\xi, 3})) \quad (3.12)$$

Proof. By assumption on the initial data, there exists $q > 3$ such that $f_0 \in L_{x,v}^q$; thanks to the a priori L^q estimate, we get $g_\epsilon \in L_{t, \text{loc}}^\infty(L_{x,v}^q)$. Define γ by:

$$\frac{1}{\gamma} = \frac{2}{3} + \frac{1}{q}$$

The previous lemma shows that for some $s > 0$:

$$\eta_\epsilon \in W_{t, \text{loc}}^{s,\gamma}(W_{x_\parallel, \text{loc}}^{s,\gamma}(W_{x_\perp, v_\perp, \text{loc}}^{-1,\gamma})) \text{ uniformly in } \epsilon$$

Since $g_\epsilon \in L_{t, \text{loc}}^q(L_{x,v}^q)$ and Ψ has compact support, we get by Hölder's inequality:

$$\eta_\epsilon \in L_{t, \text{loc}}^q(L_{x_\parallel}^q(L_{x_\perp, v_\perp}^q))$$

Since $\frac{1}{\gamma} > \frac{2}{3} > \frac{1}{3}$ and $\frac{1}{q} < \frac{1}{3}$, there exists $\theta \in]0, 1[$ such that

$$\frac{1}{3} = \frac{1-\theta}{q} + \frac{\theta}{\gamma}$$

By interpolation ([4]) we deduce that:

$$\eta_\epsilon \in W_{t, \text{loc}}^{s\theta, 3}(W_{x_\parallel, \text{loc}}^{s\theta, 3}(W_{x_\perp, v_\perp, \text{loc}}^{-\theta, 3}))$$

We then use the following refined interpolation result proved by Simon in [23], which is, roughly speaking, an anisotropic adaptation of the classical Riesz-Fréchet-Kolmogorov criterion for compactness in L^p :

Theorem 3.2. *Let $1 \leq p \leq \infty$ and $s > 0$. Let $T > 0$ and B_1, B_2 be two Banach Spaces with B_1 compactly embedded in B_2 . Let F be a bounded set of $W_t^{s,p}([0, T], B_1)$.*

Then F is relatively compact in $L_t^p([0, T], B_2)$.

This entails, thanks to Sobolev's embeddings, that the sequence (η_ϵ) is relatively compact in $L_{t, \text{loc}}^3(L_{x_\parallel, \text{loc}}^3(W_{x_\perp, v_\perp, \text{loc}}^{-\theta-\xi, 3}))$, for all $\xi > 0$. □

From now on, we consider ξ such that $\theta + \xi < 1$, which is of course possible since $\theta < 1$.

Remark. Following the remark in Step 2 and by uniqueness of the limit in the sense of distributions, we get:

$$\eta = \int G\Psi(v_{\parallel})dv_{\parallel}$$

Step 4: Passing to the weak limit

We will first need a technical lemma which is obtained directly from the 2-scale convergence of E_{ϵ} .

Lemma 3.1. *Up to a subsequence,*

- $\mathcal{R}(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v)$ 2-scale converges to $\frac{1}{2\pi} \int_0^{2\pi} \mathcal{R}(\tau)\mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v)d\tau \in L_t^{\infty}(L_{2\pi, \tau}^{\infty}(L_{x_{\parallel}}^{3/2}(W_{x_{\perp}}^{1,3/2})))$
- $R(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v)$ 2-scale converges to $\frac{1}{2\pi} \int_0^{2\pi} R(\tau)\mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v)d\tau \in L_t^{\infty}(L_{2\pi, \tau}^{\infty}(L_{x_{\parallel}}^{3/2}(W_{x_{\perp}}^{1,3/2})))$

Proof. E_{ϵ} is uniformly bounded in $L_t^{\infty}(L_{x_{\parallel}}^{3/2}(W_{x_{\perp}}^{1,3/2}))$, so there exists $\mathcal{E} \in L_t^{\infty}(L_{x_{\parallel}}^{3/2}(W_{x_{\perp}}^{1,3/2}))$ such that E_{ϵ} 2 scale converge to \mathcal{E} .

We take $\Psi(t, \tau, x)$ a 2π -periodic w.r.t. τ test function and use the 2 scale convergence of E_{ϵ} :

$$\begin{aligned} & \int \mathcal{R}(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \Psi(t, t/\epsilon, x) dt dx \\ &= \int E_{\epsilon}(t, x) \cdot {}^t\mathcal{R}(t/\epsilon)\Psi(t, t/\epsilon, x - \mathcal{R}(-t/\epsilon)v) dt dx \\ &\rightarrow \frac{1}{2\pi} \int \int_0^{2\pi} \mathcal{E}(t, \tau, x) \cdot {}^t\mathcal{R}(\tau)\Psi(t, \tau, x - \mathcal{R}(-\tau)v) dt d\tau dx \\ &= \frac{1}{2\pi} \int \int_0^{2\pi} \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) \cdot {}^t\mathcal{R}(\tau)\Psi(t, \tau, x) dt d\tau dx \end{aligned}$$

The proof is the same for $R(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v)$. □

Now, we can write the weak formulation of the kinetic equation (3.6) against a smooth test function of the form $\Phi(t, x, v_{\perp})\Psi(v_{\parallel})$ with compact support. If we can pass to the limit for such test functions, then by density it will be also the case for all test functions.

Noticing that $\operatorname{div}_x v_{\parallel} = 0$ and that

$$\operatorname{div}_x (\mathcal{R}(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v)) + \operatorname{div}_v (R(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v)) = 0$$

we get:

$$\begin{aligned} & \int \left(\partial_t (\Phi(t, x, v_{\perp})\Psi(v_{\parallel})) + v_{\parallel} \cdot \nabla_x (\Phi\Psi) + \mathcal{R}(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x (\Phi\Psi) \right. \\ & \quad \left. + R(t/\epsilon)E_{\epsilon}(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v (\Phi\Psi) \right) g_{\epsilon} dt dx_{\perp} dx_{\parallel} dv_{\perp} dv_{\parallel} \\ &= - \int u_0 \Phi(0, x, v_{\perp}) \Psi(v_{\parallel}) dx dv \end{aligned}$$

We can easily take weak limits in the linear part $\partial_t g_\epsilon + v_\parallel \cdot \nabla_x g_\epsilon$.
Consider now the “non linear” term:

$$\begin{aligned} & \int \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x \Phi(t, x, v_\perp) g_\epsilon \Psi(v_\parallel) dt dx_\perp dx_\parallel dv_\perp dv_\parallel = \\ & \int \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x \Phi(t, x, v_\perp) \left(\int g_\epsilon \Psi(v_\parallel) dv_\parallel \right) dt dx_\perp dx_\parallel dv_\perp \end{aligned}$$

The convergence of this term can be established by the strong/weak convergence principle. Nevertheless, we have to carefully use this technique to get the result and we will explicitly evaluate the difference:

$$\begin{aligned} & \left| \int \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x \Phi(t, x, v_\perp) \eta_\epsilon dt dx_\perp dx_\parallel dv_\perp \right. \\ & \quad \left. - \int \frac{1}{2\pi} \int_0^{2\pi} \mathcal{R}(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \cdot \nabla_x \Phi \eta dt dx dv_\perp \right| \\ & \leq \left| \int \left(\mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) - \frac{1}{2\pi} \int_0^{2\pi} \mathcal{R}(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_x \Phi \eta dt dx dv_\perp \right| \\ & \quad + \left| \int \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x \Phi(t, x, v_\perp) (\eta_\epsilon - \eta) dt dx_\perp dx_\parallel dv_\perp \right| \end{aligned}$$

The first term of the right hand side converges to zero because of the 2-scale convergence of E_ϵ (Lemma 3.1). We can control the second term by:

$$C \|E_\epsilon \cdot \nabla_x \Phi\|_{L_t^{3/2}(L_{x_\parallel}^{3/2}(W_{x_\perp, v_\perp}^{\theta+\xi, 3/2}))} \|\eta_\epsilon - \eta\|_{L_t^3(L_{x_\parallel}^3(W_{x_\perp, v_\perp}^{-\theta-\xi, 3}))} \quad (3.13)$$

(these norms are actually taken on the compact support of Φ but we do not write it for the sake of simplicity)

Using the fact that E_ϵ is uniformly bounded in $L_{t, \text{loc}}^{3/2}(L_{x_\parallel}^{3/2}(W_{x_\perp}^{\theta+\xi, 3/2}))$ (this is an easy consequence of Lemma 2.4) and Corollary 3.1,

$$\|\eta_\epsilon - \eta\|_{L_t^3([0, T], L_{x_\parallel}^3(K_{x_\parallel}, W_{x_\perp, v_\perp}^{-\theta-\xi, 3}(K_{x_\perp} \times K_{v_\perp})))} \rightarrow 0,$$

we can deduce that

$$\left| \int \mathcal{R}(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_x \Phi(t, x, v_\perp) (\eta_\epsilon - \eta) dt dx_\perp dx_\parallel dv_\perp \right| \rightarrow 0$$

The proof is of course the same for the other non-linear term:

$$R(t/\epsilon) E_\epsilon(t, x + \mathcal{R}(-t/\epsilon)v) \cdot \nabla_v g_\epsilon$$

To conclude let us compute the asymptotic equation satisfied by the 2-scale limit of V_ϵ denoted by V . We take $\Psi(t, \tau, x)$ a 2π -periodic w.r.t. τ test function. We write the weak formulation of the Poisson equation:

$$\begin{aligned} & \int V_\epsilon \nabla_{x_\parallel} \Psi(t, t/\epsilon, x) dt dx + \\ & \quad \epsilon^2 \int \nabla_{x_\parallel} V_\epsilon \nabla_{x_\parallel} \Psi(t, t/\epsilon, x) dt dx + \int \nabla_{x_\perp} V_\epsilon \nabla_{x_\perp} \Psi(t, t/\epsilon, x) dt dx \\ & = \int f_\epsilon(t, x, v) \Psi(t, t/\epsilon, x) dt dv dx - \int \left(\int f_0 dv dx \right) \Psi(t, t/\epsilon, x) dt dv dx \end{aligned}$$

We then pass to the 2 scale limit:

$$\begin{aligned}
& \frac{1}{2\pi} \int \int_0^{2\pi} V(t, \tau, x) \nabla_{x_\perp} \Psi(t, \tau, x) d\tau dt dx + 0 \\
+ & \frac{1}{2\pi} \int \int_0^{2\pi} \nabla_{x_\perp} V(t, \tau, x) \nabla_{x_\perp} \Psi(t, \tau, x) d\tau dt dx = \frac{1}{2\pi} \int \int_0^{2\pi} F(t, \tau, x, v) \Psi(t, \tau, x) dt dv dx \\
- & \frac{1}{2\pi} \int \int_0^{2\pi} \left(\int f_0 dv dx \right) \Psi(t, \tau, x) d\tau dv dx \\
= & \frac{1}{2\pi} \int \int_0^{2\pi} G(t, \tau, x + \mathcal{R}(\tau)v, R(\tau)v) \Psi(t, \tau, x) d\tau dv dx \\
- & \frac{1}{2\pi} \int \int_0^{2\pi} \left(\int f_0 dv dx \right) \Psi(t, \tau, x) d\tau dv dx
\end{aligned}$$

from which we get the ‘‘Poisson’’ equation given in Theorem 1.4:

$$V - \Delta_\perp V = \int G(t, x + \mathcal{R}(\tau)v, R(\tau)v) dv - \int f_0 dv dx$$

Moreover since $E_{\epsilon, \perp} = -\nabla_{x_\perp} V_\epsilon$ and thanks to Remark 1 following Lemma 2.4, we easily get if we pass to the two-scale limit:

$$\mathcal{E} = (-\nabla_{x_\perp} V, 0)$$

□

4 Concluding comments

4.1 Comments on the result

Finally we can see as in [12] (namely by performing the change of variables $x = x_c - v^\perp$ and looking at the new equations in the so-called gyro-variables (t, x_c, v)) that the drift involving the electric field in the asymptotic ‘‘kinetic’’ equation corresponds to the electric drift that we mentioned in the introduction and which was expected to appear. We also notice that the Poisson equation we get in the asymptotic system is the same than the one used in the numerical simulations of tokamak plasmas (see for example the GYSELA code in [18]). Nevertheless, physicists do not get it in the same formal way: they claim that it only expresses the quasineutrality of the plasma (there is no ‘‘real’’ Poisson equation involved) and the perpendicular laplacian happens to appear due from the so-called ‘‘polarization drift’’ ([18], see also [24] for a physical reference on the subject). It would be interesting to justify such a computation from a mathematical point of view.

At last, we wish to point out a really unpleasant feature of our model, which is that there is no parallel dynamics.

4.2 An alternative model

Let us give some comments on the gyrokinetic approximation of the system (4.2) which consists in considering a fixed background of electrons:

$$n_\epsilon^e = \int f_0 dx dv \tag{4.1}$$

Actually, quite surprisingly, this model engenders more interesting physical properties: in this case the parallel component of the electric field does not vanish but appears as a pressure in the end (which may bring difficulties both in the study of the asymptotic system and in numerical simulations).

$$\begin{cases} \partial_t f_\epsilon + \frac{v_\perp}{\epsilon} \cdot \nabla_x f_\epsilon + v_\parallel \cdot \nabla_x f_\epsilon + (E_\epsilon + \frac{v \wedge B}{\epsilon}) \cdot \nabla_v f_\epsilon = 0 \\ E_\epsilon = (-\nabla_{x_\perp} V_\epsilon, -\epsilon \nabla_{x_\parallel} V_\epsilon) \\ -\epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon = \int f_\epsilon dv - \int f_0 dv dx \\ f_{\epsilon, t=0} = f_{\epsilon, 0} \end{cases} \quad (4.2)$$

With the same computations as the present paper, we get:

$$f_\epsilon \text{ 2-scale converges to } F \quad (4.3)$$

$$E_\epsilon \text{ 2-scale converges to } \mathcal{E} \quad (4.4)$$

In a formal sense, there exists a function G such that:

$$F(t, \tau, x, v) = G(t, x + \mathcal{R}(\tau)v, R(\tau)v) \quad (4.5)$$

and (G, \mathcal{E}) is solution to:

$$\begin{aligned} \partial_t G + v_\parallel \cdot \nabla_x G + \frac{1}{2\pi} \left(\int_0^{2\pi} \mathcal{R}(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_x G \\ + \frac{1}{2\pi} \left(\int_0^{2\pi} R(\tau) \mathcal{E}(t, \tau, x + \mathcal{R}(-\tau)v) d\tau \right) \cdot \nabla_v G = 0 \\ G|_{t=0} = f_0 \end{aligned}$$

$$\mathcal{E} = (-\nabla_\perp V, \mathcal{E}_\parallel), \quad -\Delta_\perp V = \int G(t, x + \mathcal{R}(\tau)v, R(\tau)v) dv - \int f_0 dv dx$$

still denoting by R and \mathcal{R} the linear operators defined by:

$$R(\tau) = \begin{bmatrix} \cos \tau & -\sin \tau & 0 \\ \sin \tau & \cos \tau & 0 \\ 0 & 0 & 1 \end{bmatrix}, \quad \mathcal{R}(\tau) = (R(-\pi/2) - R(-\pi/2 + \tau))$$

The parallel component \mathcal{E}_\parallel has to be seen as a pressure (or the Lagrange multiplier) associated to the ‘‘incompressibility’’ constraint $\int_{\mathbb{T}^2} \int G(t, x + \mathcal{R}(\tau)v, R(\tau)v) dv dx_\perp = \int f_0 dv dx$

Let us just give a few words on the difficulties that arise with this model. The Poisson equation can be restated as:

$$-\epsilon^2 \Delta_{x_\parallel} V_\epsilon - \Delta_{x_\perp} V_\epsilon = \int f_\epsilon dv - \int f_\epsilon dv dx_\perp + \int f_\epsilon dv dx_\perp - \int f_0 dv dx \quad (4.6)$$

so that thanks to the linearity of the Poisson equation we can study separately two equations. The first one states:

$$-\epsilon^2 \Delta_{x_\parallel} V_\epsilon^1 - \Delta_{x_\perp} V_\epsilon^1 = \int f_\epsilon dv - \int f_\epsilon dv dx_\perp \quad (4.7)$$

For this part of the electric potential we get the same estimates as in lemma 2.4. Indeed, $\int (\int f_\epsilon dv - \int f_\epsilon dv dx_\perp) dx_\perp = 0$ so that we can use elliptic estimates on the torus $\mathbb{T}_{x_\perp}^2$.

Consequently this electric potential does not give birth to any parallel dynamics, like in Theorem 1.4.

The second one is:

$$-\epsilon^2 \Delta_{x_{\parallel}} V_{\epsilon}^2 = \underbrace{\int f_{\epsilon} dv dx_{\perp} - \int f_0 dv dx}_{\text{only depends on } x_{\parallel}} \quad (4.8)$$

This equation, associated to the one giving the electric field $E_{\epsilon, \parallel}^2 = -\epsilon \partial_{x_{\parallel}} V_{\parallel}^2$, is similar to the one studied by Grenier in [19], coupled to a Vlasov equation describing a quasineutral plasma. In this case it was shown that there exist plasma waves with both temporal and spatial oscillations with frequency $\frac{1}{\sqrt{\epsilon}}$ and magnitude of order $\frac{1}{\sqrt{\epsilon}}$. Because of these waves, it is much more difficult to pass to the limit in order to get a kinetic equation. Grenier managed to prove the convergence only for distribution functions with special form and got in the end a Euler-like system with an electric field interpreted as a Lagrange multiplier. Hence, in our case, this part of the electric field may engender a non-zero \mathcal{E}_{\parallel} .

For these reasons, it seems much harder to expect to prove a rigorous result similar to Theorem 1.4 with such a model.

4.3 Prospects

A way to pass to the limit in this latest case would be to use a relative entropy method like in the papers of Brenier ([5]) and Golse and Saint-Raymond ([16]). This will be the object of a forthcoming work.

Another interesting issue would be to consider a “true” Boltzmann-Maxwell distribution for the electrons (not linearized like in this paper) and perform an asymptotic analysis, maybe also with a relative entropy method.

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References

- [1] V.I. Agoshkov. Spaces of functions with differential-difference characteristics and the smoothness of solutions of the transport equation. *Dokl. Akad. Nauk SSSR*, 276(6):1289–1293, 1984.
- [2] G. Allaire. Homogenization and Two-scale Convergence. *SIAM J. Math. Anal.*, XXIII(6):1482–1518, 1992.
- [3] A.A. Arsenev. Existence in the large of a weak solution of Vlasov’s system of equations. *Z. Vychisl. Mat. Mat. Fiz.*, 15:136–147, 1975.
- [4] J. Bergh and J. Löfström. Interpolation Spaces. An introduction. *Springer, Berlin*, 1976.
- [5] Y. Brenier. Convergence of the Vlasov-Poisson system to the incompressible Euler equations. *Comm. P.D.E.*, 25:737–754, 2000.
- [6] M. Bézard. Régularité L^p précisée des moyennes dans les équations de transport. *Bull. Soc. math. France*, 122:29–76, 1994.

- [7] P. Degond. Global existence of smooth solutions for the Vlasov-Fokker-Planck equations in 1 and 2 space dimensions. *Ann. Sci. Ecole Norm. Sup.*, 19:519–542, 1986.
- [8] R.J DiPerna and P.L. Lions. Global weak solutions of Vlasov-Maxwell systems. *Comm. Pure App. Math.*, 42:729–757, 1989.
- [9] R.J. DiPerna, P.L. Lions, and Y. Meyer. L^p regularity of velocity averages. *Annales de l'I.H.P, section C*, 8(3-4):271–287, 1991.
- [10] E. Frénod, P.A. Raviart, and E. Sonnendrücker. Asymptotic expansion of the Vlasov equation in a large external magnetic field. *J. Math. Pures et Appl.*, 80(8):2001, 815-843.
- [11] E. Frénod and E. Sonnendrücker. Homogenization of the Vlasov equation and of the Vlasov-Poisson system with a strong external magnetic field. *Asymp. Anal.*, 18(3-4):193–214, 1998.
- [12] E. Frénod and E. Sonnendrücker. The Finite Larmor Radius Approximation. *SIAM J. Math. Anal.*, 32(6):1227–1247, 2001.
- [13] F. Golse, P.L. Lions, B. Perthame, and R. Sentis. Regularity of the moments of the solution of a transport equation. *J. Funct. Anal.*, 26:110–125, 1988.
- [14] F. Golse, B. Perthame, and R. Sentis. Un résultat de compacité pour les equations de transport et application au calcul de la limite de la valeur propre principale d'un opérateur de transport. *C.R. Acad. Sci. Paris Sér. I Math.*, 301(7):341–344, 1985.
- [15] F. Golse and L. Saint-Raymond. The Vlasov-Poisson system with strong magnetic field. *J. Math. Pures. Appl.*, 78:791–817, 1999.
- [16] F. Golse and L. Saint-Raymond. The Vlasov-Poisson system with strong magnetic field in quasineutral regime. *Mathematical Models and Methods in Applied Sciences*, 13(5):661–714, 2003.
- [17] H. Grad. The guiding center plasma. *Proc. Sympos. Appl. Math. Amer. Math. Soc.*, 18:162–248, 1967.
- [18] V. Grandgirard et al. Global full- f gyrokinetic simulations of plasma turbulence. *Plasma Phys. Control. Fusion*, 49:173–182, 2007.
- [19] E. Grenier. Oscillations in quasineutral plasmas. *Commun. in partial differential equations*, 21(3-4):363–394, 1996.
- [20] G. Nguetseng. A general convergence result for a functional related to the theory of homogenization. *SIAM J. Math. Anal.*, 20(3):608–623, 1989.
- [21] L. Saint-Raymond. The gyrokinetic approximation for the Vlasov-Poisson system. *Math. Mod. Meth. Appl. Sci.*, 10(9):1305–1332, 2000.
- [22] L. Saint-Raymond. Control of large velocities in the two-dimensional gyrokinetic approximation. *J. Math. Pures Appl.*, 81:379–399, 2002.
- [23] J. Simon. Compact sets in $L^p(0, T; B)$. *Ann. Mat. Pura. Appl. IV*, 146:65–96, 1987.
- [24] J. Wesson. Tokamaks. *Clarendon Press-Oxford*, 2004.