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***The Vlasov-Maxwell system with strong initial
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The Vlasov-Maxwell system with strong initial magnetic field. Part I : Guiding-center approximation

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Abstract: In this paper we study the asymptotic behavior of the Vlasov-Maxwell equations with strong magnetic field. More precisely we investigate the Cauchy problems associated to strong initial magnetic fields. We justify the convergence towards the so-called "guiding center approximation" when the dynamics is observed on a slower time scale than the plasma frequency. Our proofs rely on the modulated energy method.

Key-words: Vlasov-Maxwell equations, Guiding-center approximation, Modulated energy

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Equations de Vlasov-Maxwell avec champ magnétique initial fort. Partie I : approximation de centre-guide

Résumé : Dans cet article nous étudions le comportement asymptotique des équations de Vlasov-Maxwell avec champ magnétique fort. Plus précisément nous analysons les problèmes de Cauchy lorsque le champ magnétique initial tend vers l'infini. On justifie la convergence vers l'approximation "centre-guide" quand l'échelle d'observation en temps est petite devant la fréquence des oscillations du plasma. Les preuves reposent sur la méthode de l'énergie modulée.

Mots-clés : Equations de Vlasov-Maxwell, Approximation centre-guide, Energie modulée

1 Introduction

The main motivations and applications in plasma physics concern the energy production through the thermonuclear fusion process. Two ways are currently explored for this: the inertial confinement fusion (ICF) and the magnetic confinement fusion (MCF). The magnetic confinement is performed in large toroidal devices, called tokamaks, by using strong magnetic fields. Besides studying these phenomena by direct observation and measurements, the numerical simulation of them is of crucial importance.

The dynamics of charged particles is described in terms of a number density by the Vlasov equation, coupled to the Maxwell equations for the electro-magnetic field. Generally the numerical resolution of this model requires important computational efforts, since we are working in a phase space with three spatial dimensions and three momentum dimensions. Moreover new difficulties appear when studying strong magnetic field regimes: large magnetic fields introduce a new time scale, related to the period of rotation of the particles around the magnetic field lines. Since the cyclotron period is proportional to the inverse of the magnitude of the magnetic field, the above time scale is very restrictive from the numerical point of view. Hence it is worth looking for simpler approximate models, like the gyro-kinetic model or the guiding center model [15], [19]. The limits of the Vlasov or Vlasov-Poisson equations with strong external magnetic fields have been investigated recently [7], [11], [8], [3]. For related works we refer to [16], [17].

For understanding the effects of strong magnetic fields let us start by analyzing the motion of individual charged particles under the action of constant electro-magnetic field (E, B) . The motion equations of a particle of mass m and charge q are given by

$$\frac{dX}{ds} = V(s), \quad \frac{dV}{ds} = \frac{q}{m}(E + V(s) \wedge B),$$

where $(X(s), V(s))$ represent the position and velocity at time s . Projecting on the direction of B it is easily seen that

$$\frac{d}{ds} \left(V(s) \cdot \frac{B}{|B|} \right) = \frac{q}{m} \frac{E \cdot B}{|B|},$$

saying that the particle is advected with the acceleration $\frac{q}{m} \frac{E \cdot B}{|B|}$ in the direction of B . Note that this acceleration do not depend on the magnitude of B . For analyzing the motion in the plane orthogonal to B it is convenient to represent the velocity as $V(s) = \frac{E \wedge B}{|B|^2} + U(s)$ where U satisfies

$$\frac{dU}{ds} = \frac{q}{m} \left((E \cdot B) \frac{B}{|B|^2} + U(s) \wedge B \right).$$

We denote by U_{\perp} the projection of U on the plane orthogonal to B

$$U_{\perp}(s) = \left(\frac{B}{|B|} \wedge U(s) \right) \wedge \frac{B}{|B|}.$$

A straightforward computation shows that U_\perp verifies

$$\frac{d^2}{ds^2}U_\perp + \frac{q^2}{m^2}|B|^2U_\perp(s) = 0,$$

implying that

$$U_\perp(s) = \mathcal{R}(-\omega_c s)U_\perp(0) = \mathcal{R}(-\omega_c s) \left(V_\perp(0) - \frac{E \wedge B}{|B|^2} \right),$$

where $\omega_c = \frac{|q|}{m}|B|$ is the cyclotron frequency and for any $\theta \in \mathbb{R}$ we denote by $\mathcal{R}(\theta)$ the rotation of angle θ in the plane orthogonal to B , oriented by qB . We deduce that

$$X_\perp(t) = X_\perp(0) - \frac{1}{\omega_c} \mathcal{R} \left(\frac{\pi}{2} \right) U_\perp(0) + t \frac{E \wedge B}{|B|^2} + \frac{1}{\omega_c} \mathcal{R} \left(-\omega_c t + \frac{\pi}{2} \right) U_\perp(0).$$

The particles move on a helix with axis parallel to B and radius (called the Larmor radius) proportional to $\frac{1}{\omega_c} = \frac{m}{|q||B|}$. Therefore, when the magnetic field is large, the Larmor radius goes to zero and the particle motion can be approximated by the motion of the axis, whose velocity in the plane orthogonal to B , given by $\frac{E \wedge B}{|B|^2}$, is called the drift velocity. Notice that the drift velocity associated to strong magnetic fields $B = \mathcal{O}(1/\varepsilon)$ is small $\frac{E \wedge B}{|B|^2} = \mathcal{O}(\varepsilon)$. Hence the motion of the axis becomes significant only for large observation time $\mathcal{O}(1/\varepsilon)$.

We consider a population of relativistic electrons whose density in the phase space is denoted by f . We neglect the collisions between particles assuming that they interact only by electro-magnetic fields created collectively. The particle density depends on time $t \in \mathbb{R}_+$, position $x \in \mathbb{R}^3$, momentum $p \in \mathbb{R}^3$ and satisfies the Vlasov equation

$$\partial_t f + v(p) \cdot \nabla_x f - e(E(t, x) + v(p) \wedge B(t, x)) \cdot \nabla_p f = 0, \quad (1)$$

where $-e < 0$ is the electron charge, $v(p)$ is the relativistic velocity associated to the momentum p

$$v(p) = \frac{p}{m_e} \left(1 + \frac{|p|^2}{m_e^2 c_0^2} \right)^{-\frac{1}{2}},$$

m_e is the electron mass and c_0 is the vacuum light speed. The electro-magnetic field is defined in a self-consistent way by the Maxwell equations

$$\partial_t E - c_0^2 \operatorname{curl}_x B = \frac{e}{\varepsilon_0} \int_{\mathbb{R}^3} v(p) f(t, x, p) dp, \quad (2)$$

$$\partial_t B + \operatorname{curl}_x E = 0, \quad (3)$$

$$\operatorname{div}_x E = \frac{e}{\varepsilon_0} \left(n - \int_{\mathbb{R}^3} f(t, x, p) dp \right), \quad \operatorname{div}_x B = 0, \quad (4)$$

where ε_0 is the vacuum permittivity and n is the concentration of a background ion distribution (*i.e.*, the number of ions per volume unit). Let us write the equations in dimensionless form. We define the thermal potential by $U_{\text{th}} = \frac{K_B T_{\text{th}}}{e}$ where K_B is the Boltzmann constant and T_{th} is the temperature. The thermal momentum p_{th} is given by

$$m_e c_0^2 \left(\left(1 + \frac{p_{\text{th}}^2}{m_e^2 c_0^2} \right)^{\frac{1}{2}} - 1 \right) = K_B T_{\text{th}},$$

which leads to

$$p_{\text{th}} = \left((K_B T_{\text{th}})^2 / c_0^2 + 2K_B T_{\text{th}} m_e \right)^{\frac{1}{2}}.$$

We introduce a length unit L and a time unit T . As momentum unit we set $P = p_{\text{th}}$. We define dimensionless variables and unknowns by the relations

$$t = Tt', \quad x = Lx', \quad p = p_{\text{th}}p',$$

$$f(t, x, p) = \frac{n_e}{p_{\text{th}}^3} f' \left(\frac{t}{T}, \frac{x}{L}, \frac{p}{p_{\text{th}}} \right), \quad E(t, x) = \frac{U_{\text{th}}}{L} E' \left(\frac{t}{T}, \frac{x}{L} \right), \quad B(t, x) = \frac{1}{\varepsilon} \frac{m_e}{e T_p} B' \left(\frac{t}{T}, \frac{x}{L} \right),$$

where $T_p = \left(\frac{m_e \varepsilon_0}{e^2 n_e} \right)^{1/2}$ is the inverse of the plasma frequency, n_e is the average of the electron concentration and $\varepsilon > 0$ is a small parameter. We assume that the plasma is globally neutral and therefore we have $n_e = n$. We set

$$v'(p') = \frac{p_{\text{th}}^2}{m_e K_B T_{\text{th}}} p' \left(1 + \frac{p_{\text{th}}^2}{m_e^2 c_0^2} |p'|^2 \right)^{-\frac{1}{2}}.$$

As a matter of fact note that $v(p) = \frac{K_B T_{\text{th}}}{p_{\text{th}}} v'(p/p_{\text{th}})$. We also introduce the Debye length

$$\lambda_D = \left(\frac{\varepsilon_0 K_B T_{\text{th}}}{e^2 n} \right)^{\frac{1}{2}}.$$

Notice that we have $K_B T_{\text{th}} / m_e = (\lambda_D / T_p)^2$. Then the equations become having dropped the primes

$$\partial_t f + \frac{K_B T_{\text{th}}}{p_{\text{th}}} \frac{T}{L} v(p) \cdot \nabla_x f - \frac{K_B T_{\text{th}}}{p_{\text{th}}} \frac{T}{L} \left(E(t, x) + \frac{L m_e}{T_p p_{\text{th}}} v(p) \wedge \frac{B(t, x)}{\varepsilon} \right) \cdot \nabla_p f = 0,$$

$$\partial_t E - \frac{T}{T_p} \frac{m_e c_0^2}{K_B T_{\text{th}}} \text{curl}_x \left(\frac{B}{\varepsilon} \right) = \left(\frac{L}{\lambda_D} \right)^2 \frac{K_B T_{\text{th}}}{p_{\text{th}}} \frac{T}{L} j(t, x),$$

$$\partial_t \left(\frac{B}{\varepsilon} \right) + \frac{T}{T_p} \left(\frac{\lambda_D}{L} \right)^2 \text{curl}_x E = 0,$$

$$\operatorname{div}_x E = \left(\frac{L}{\lambda_D} \right)^2 (1 - \rho(t, x)), \quad \operatorname{div}_x B = 0,$$

where $\rho(t, x) = \int_{\mathbb{R}^3} f(t, x, p) dp$, $j(t, x) = \int_{\mathbb{R}^3} v(p) f(t, x, p) dp$. We take as length unit $L = \lambda_D$ and as time unit $T = \frac{T_p}{\varepsilon}$. Observe that

$$\frac{K_B T_{\text{th}}}{p_{\text{th}}} \frac{T_p}{\lambda_D} = \frac{\lambda_D m_e}{T_p p_{\text{th}}} = \left(\frac{K_B T_{\text{th}}}{m_e c_0^2} + 2 \right)^{-\frac{1}{2}} =: \alpha.$$

Finally we obtain the equations

$$\partial_t f + \frac{\alpha}{\varepsilon} v(p) \cdot \nabla_x f - \frac{\alpha}{\varepsilon} \left(E(t, x) + \alpha v(p) \wedge \frac{B(t, x)}{\varepsilon} \right) \cdot \nabla_p f = 0, \quad (5)$$

$$\partial_t E - \frac{1}{\varepsilon \beta} \operatorname{curl}_x \left(\frac{B}{\varepsilon} \right) = \frac{\alpha}{\varepsilon} j(t, x), \quad (6)$$

$$\partial_t \left(\frac{B}{\varepsilon} \right) + \frac{1}{\varepsilon} \operatorname{curl}_x E = 0, \quad (7)$$

$$\operatorname{div}_x E = 1 - \rho(t, x), \quad \operatorname{div}_x B = 0, \quad (8)$$

with $\beta = \frac{K_B T_{\text{th}}}{m_e c_0^2}$ and $v(p) = \frac{p}{\alpha^2} \left(1 + \frac{\beta}{\alpha^2} |p|^2 \right)^{-1/2}$. We are concerned with the asymptotic behavior of (5), (6), (7), (8) when $\varepsilon \searrow 0$, $\beta = \mathcal{O}(1)$ and therefore $\alpha = \mathcal{O}(1)$. In order to simplify our computations we will study the systems

$$\partial_t f^\varepsilon + \frac{1}{\varepsilon} v(p) \cdot \nabla_x f^\varepsilon - \frac{1}{\varepsilon} \left(E^\varepsilon(t, x) + v(p) \wedge \frac{B^\varepsilon(t, x)}{\varepsilon} \right) \cdot \nabla_p f^\varepsilon = 0, \quad (9)$$

$$\partial_t E^\varepsilon - \frac{1}{\varepsilon} \operatorname{curl}_x \left(\frac{B^\varepsilon}{\varepsilon} \right) = \frac{1}{\varepsilon} j^\varepsilon(t, x), \quad (10)$$

$$\partial_t \left(\frac{B^\varepsilon}{\varepsilon} \right) + \frac{1}{\varepsilon} \operatorname{curl}_x E^\varepsilon = 0, \quad (11)$$

$$\operatorname{div}_x E^\varepsilon = 1 - \rho^\varepsilon(t, x), \quad \operatorname{div}_x B^\varepsilon = 0, \quad (12)$$

$$\rho^\varepsilon = \int_{\mathbb{R}^3} f^\varepsilon dp, \quad j^\varepsilon = \int_{\mathbb{R}^3} v(p) f^\varepsilon dp, \quad v(p) = \frac{p}{(1 + |p|^2)^{\frac{1}{2}}}, \quad (13)$$

which has the same structure as (5), (6), (7), (8). We prescribe also the initial conditions

$$f^\varepsilon(0, x, p) = f_0^\varepsilon(x, p), \quad (E^\varepsilon, B^\varepsilon)(0, x) = (E_0^\varepsilon, B_0^\varepsilon)(x). \quad (14)$$

We assume also periodicity in the space variable $x \in \mathbb{T}^d$ where $\mathbb{T}^d = \mathbb{R}^d / \mathbb{Z}^d$, equipped with the restriction of the Lebesgue measure of \mathbb{R}^d on $[0, 1]^d$, $d \in \{1, 2, 3\}$. The subject matter of this paper concerns the stability of the solutions $(f^\varepsilon, E^\varepsilon, B^\varepsilon)_{\varepsilon>0}$ for well prepared initial conditions $(f_0^\varepsilon, E_0^\varepsilon, B_0^\varepsilon)_{\varepsilon>0}$, where $\varepsilon > 0$ is a small parameter. Consider a constant magnetic field $B_0 \in \mathbb{R}^3$ and observe that (10), (11) can be written in the form

$$\partial_t E^\varepsilon - \frac{1}{\varepsilon} \operatorname{curl}_x \left(\frac{B^\varepsilon - B_0}{\varepsilon} \right) = \frac{j^\varepsilon(t, x)}{\varepsilon}, \quad (15)$$

$$\partial_t \left(\frac{B^\varepsilon - B_0}{\varepsilon} \right) + \frac{1}{\varepsilon} \operatorname{curl}_x E^\varepsilon = 0. \quad (16)$$

Multiplying (9) by $(1 + |p|^2)^{\frac{1}{2}} - 1$, (15) by E^ε and (16) by $\left(\frac{B^\varepsilon - B_0}{\varepsilon} \right)$ one gets as usual the conservation

$$\frac{d}{dt} \left\{ \int_{\mathbb{T}^3} \int_{\mathbb{R}^3} ((1 + |p|^2)^{\frac{1}{2}} - 1) f^\varepsilon dp dx + \frac{1}{2} \int_{\mathbb{T}^3} \left(|E^\varepsilon|^2 + \left| \frac{B^\varepsilon - B_0}{\varepsilon} \right|^2 \right) dx \right\} = 0, \quad (17)$$

implying that

$$\begin{aligned} \int_{\mathbb{T}^3} \left| \frac{B^\varepsilon(t, x) - B_0}{\varepsilon} \right|^2 dx &\leq 2 \int_{\mathbb{T}^3} \int_{\mathbb{R}^3} ((1 + |p|^2)^{\frac{1}{2}} - 1) f_0^\varepsilon(x, p) dp dx \\ &\quad + \int_{\mathbb{T}^3} |E_0^\varepsilon(x)|^2 dx + \int_{\mathbb{T}^3} \left| \frac{B_0^\varepsilon(x) - B_0}{\varepsilon} \right|^2 dx. \end{aligned}$$

In particular we deduce that $\sup_{\varepsilon>0, t \in \mathbb{R}_+} \int_{\mathbb{T}^3} \varepsilon^{-2} |B^\varepsilon(t, x) - B_0|^2 dx < +\infty$ for initial conditions satisfying

$$\sup_{\varepsilon>0} \left\{ \int_{\mathbb{T}^3} \int_{\mathbb{R}^3} ((1 + |p|^2)^{\frac{1}{2}} - 1) f_0^\varepsilon dp dx + \frac{1}{2} \int_{\mathbb{T}^3} |E_0^\varepsilon|^2 dx + \frac{1}{2} \int_{\mathbb{T}^3} \left| \frac{B_0^\varepsilon - B_0}{\varepsilon} \right|^2 dx \right\} < +\infty.$$

Recalling that the unit for the magnetic field was chosen proportional to $1/\varepsilon$, the above arguments say that if initially the (unscaled) magnetic field is close to $\frac{B_0}{\varepsilon}$, then at any time $t > 0$ the (unscaled) magnetic field remains close to $\frac{B_0}{\varepsilon}$; we are dealing with a strong magnetic field regime. As a matter of fact this regime is consistent with the Vlasov-Poisson equations with strong external magnetic field. This asymptotic regime has been investigated in [11] by appealing to compactness methods. In the two dimensional case the authors justified the convergence towards the vorticity formulation of the incompressible Euler equations with a right-hand side involving a defect measure. Another approach uses modulated energy (or relative entropy) methods, as introduced in [24]. By this technique one gets strong convergences, provided that the solution of the limit system is smooth. Results for the Vlasov-Poisson equations with strong magnetic field have been obtained recently in [3], [12]. More generally the relative entropy method allows the treatment of various asymptotic

questions in plasma physics [4], [14], [2], gas dynamics [22], [1], fluid-particles interaction [13].

We intend to address the Vlasov-Maxwell system with strong initial magnetic field by the method of relative entropy. We follow the ideas in [3] by adapting the arguments to the relativistic case with self-consistent magnetic field. This generalization is important from the physical point of view since we are dealing with a more realistic model. Besides, this work shows how robust the relative entropy method is, which is interesting from the mathematical point of view. A complete convergence result is obtained in the two dimensional case, see Theorem 2.1.

The paper is organized as follows. The relativistic Vlasov-Maxwell system in two dimensions is treated in Section 2. After a formal derivation of the limit system we introduce the modulated energy. We study the time evolution of it and we deduce strong convergence for the electro-magnetic field. We obtain also convergence in the distribution sense for the macroscopic quantities like the charge and current densities. The last section is devoted to other systems, as the non relativistic case or cases with particle densities depending on macroscopic charge densities and bulk velocities.

2 The two dimensional case

We consider the Vlasov-Maxwell system (9), (10), (11), (12) in two dimensions. For any $\varepsilon > 0$ we are looking for a solution with particle density $f^\varepsilon = f^\varepsilon(t, x, p)$, $(t, x, p) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2$ and electro-magnetic field of the form $((E_1^\varepsilon, E_2^\varepsilon, 0), (0, 0, B_3^\varepsilon))$. It is convenient to introduce the new momentum variable $u = \frac{p}{\varepsilon}$ and the new density function

$$F^\varepsilon(t, x, u) = \varepsilon^2 f^\varepsilon(t, x, \varepsilon u), \quad (t, x, u) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2.$$

Observe that these distributions have the same charge densities

$$\rho^\varepsilon(t, x) = \int_{\mathbb{R}^2} f^\varepsilon(t, x, p) dp = \int_{\mathbb{R}^2} F^\varepsilon(t, x, u) du,$$

and that the current densities are related by

$$j^\varepsilon(t, x) = \int_{\mathbb{R}^2} v(p) f^\varepsilon(t, x, p) dp = \varepsilon \int_{\mathbb{R}^2} v^\varepsilon(u) F^\varepsilon(t, x, u) du = \varepsilon J^\varepsilon(t, x),$$

where the velocity v^ε is given by $v^\varepsilon(u) = u/(1+\varepsilon^2|u|^2)^{\frac{1}{2}}$. We use the notation ${}^\perp v = (v_2, -v_1)$, $\forall v = (v_1, v_2) \in \mathbb{R}^2$. With these notations the two dimensional Vlasov-Maxwell system becomes

$$\partial_t F^\varepsilon + v^\varepsilon(u) \cdot \nabla_x F^\varepsilon - \frac{1}{\varepsilon^2} (E^\varepsilon(t, x) + B_3^\varepsilon(t, x) {}^\perp v^\varepsilon(u)) \cdot \nabla_u F^\varepsilon = 0, \quad (t, x, u) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2, \quad (18)$$

$$\partial_t E_1^\varepsilon - \frac{1}{\varepsilon} \partial_{x_2} \left(\frac{B_3^\varepsilon}{\varepsilon} \right) = J_1^\varepsilon(t, x), \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (19)$$

$$\partial_t E_2^\varepsilon + \frac{1}{\varepsilon} \partial_{x_1} \left(\frac{B_3^\varepsilon}{\varepsilon} \right) = J_2^\varepsilon(t, x), \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (20)$$

$$\partial_t \left(\frac{B_3^\varepsilon}{\varepsilon} \right) + \frac{1}{\varepsilon} (\partial_{x_1} E_2^\varepsilon - \partial_{x_2} E_1^\varepsilon) = 0, \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (21)$$

$$\partial_{x_1} E_1^\varepsilon + \partial_{x_2} E_2^\varepsilon = 1 - \rho^\varepsilon(t, x), \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (22)$$

with the initial conditions

$$F^\varepsilon(0, x, u) = \varepsilon^2 f_0^\varepsilon(x, \varepsilon u) =: F_0^\varepsilon(x, u), \quad (t, x, u) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2, \quad (23)$$

$$(E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)(0, x) = (E_{0,1}^\varepsilon, E_{0,2}^\varepsilon, B_{0,3}^\varepsilon)(x), \quad x \in \mathbb{T}^2. \quad (24)$$

We make the following hypotheses on the initial conditions $(f_0^\varepsilon, E_{0,1}^\varepsilon, E_{0,2}^\varepsilon, B_{0,3}^\varepsilon)$

H1) $f_0^\varepsilon \geq 0$, $\int_{\mathbb{T}^2} \int_{\mathbb{R}^2} f_0^\varepsilon(x, p) dp dx = 1$;

H2) $\lim_{\varepsilon \searrow 0} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} ((1 + |p|^2)^{\frac{1}{2}} - 1) f_0^\varepsilon(x, p) dp dx = 0$;

H3) $(E_{0,1}^\varepsilon, E_{0,2}^\varepsilon, B_{0,3}^\varepsilon) \in L^2(\mathbb{T}^2)^3$ and $\operatorname{div}_x E_0^\varepsilon = 1 - \rho_0^\varepsilon$ where $\rho_0^\varepsilon = \int_{\mathbb{R}^2} f_0^\varepsilon dp$;

H4) there are $E_0 = (E_{0,1}, E_{0,2}) \in L^2(\mathbb{T}^2)^2$ verifying $\partial_{x_1} E_{0,2} - \partial_{x_2} E_{0,1} = 0$ and a constant magnetic field $(0, 0, B_{0,3})$ with $B_{0,3} \neq 0$ such that

$$\lim_{\varepsilon \searrow 0} \left\{ \frac{1}{2} \int_{\mathbb{T}^2} |E_0^\varepsilon(x) - E_0(x)|^2 dx + \frac{1}{2} \int_{\mathbb{T}^2} \left(\frac{B_{0,3}^\varepsilon(x) - B_{0,3}}{\varepsilon} \right)^2 dx \right\} = 0.$$

Since $\operatorname{div}_x E_0^\varepsilon = 1 - \rho_0^\varepsilon$ and $\lim_{\varepsilon \searrow 0} E_0^\varepsilon = E_0$ in $L^2(\mathbb{T}^2)^2$ we deduce that $\lim_{\varepsilon \searrow 0} \rho_0^\varepsilon = 1 - \operatorname{div}_x E_0$ in $\mathcal{D}'(\mathbb{T}^2)$ and therefore the electric field $E_0 \in L^2(\mathbb{T}^2)^2$ in H4 solves the problem

$$\operatorname{div}_x^\perp E_0 = 0, \quad \operatorname{div}_x E_0 = 1 - \lim_{\varepsilon \searrow 0} \rho_0^\varepsilon \quad \text{in } \mathcal{D}'(\mathbb{T}^2).$$

Assume that the initial charge densities $(\rho_0^\varepsilon)_{\varepsilon > 0}$ are bounded in $L^r(\mathbb{T}^2)$ for some finite $r > 1$ and consider a sequence $(\varepsilon_k)_k$ converging towards zero such that $\lim_{k \rightarrow +\infty} \rho_0^{\varepsilon_k} = \rho_0$ weakly in $L^r(\mathbb{T}^2)$. In this case the electric field appearing in H4 is unique up to two constants $e_0 = (e_{0,1}, e_{0,2}) \in \mathbb{R}^2$, $E_0 = \nabla_x \phi_0 + e_0$ where $\phi_0 \in W^{2,r}(\mathbb{T}^2)$ is the unique solution of

$$-\Delta_x \phi_0 = \rho_0(x) - 1, \quad x \in \mathbb{T}^2, \quad \int_{\mathbb{T}^2} \phi_0(x) dx = 0.$$

Notice that H1, H2 are equivalent to

$$F_0^\varepsilon \geq 0, \quad \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} F_0^\varepsilon(x, u) du dx = 1, \quad \lim_{\varepsilon \searrow 0} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} ((1 + \varepsilon^2 |u|^2)^{\frac{1}{2}} - 1) F_0^\varepsilon(x, u) du dx = 0.$$

The theory for the existence and uniqueness of global classical solution for the relativistic Vlasov-Maxwell system is now well developed in two dimensions cf. [10].

2.1 Analysis of the limit system

Let $(F^\varepsilon, E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)_{\varepsilon>0}$ be smooth solutions for (18), (19), (20), (21), (22) with smooth initial conditions (23), (24). The conservation of the total energy implies

$$\begin{aligned} & \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon(t, x, u) \, du \, dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx \\ &= \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F_0^\varepsilon(x, u) \, du \, dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E_0^\varepsilon(x)|^2 + \left(\frac{B_{0,3}^\varepsilon(x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx, \end{aligned}$$

where $\mathcal{E}^\varepsilon(u) = \varepsilon^{-2}((1 + \varepsilon^2|u|^2)^{\frac{1}{2}} - 1)$ is the energy associated to the velocity $v^\varepsilon(u)$ (*i.e.*, $\nabla_u \mathcal{E}^\varepsilon = v^\varepsilon(u)$) and $B_{0,3}$ is the constant appearing in H4. We deduce that

$$\sup_{\varepsilon>0, t \in \mathbb{R}_+} \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon \, du \, dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx < +\infty. \quad (25)$$

In particular there is a sequence $(\varepsilon_k)_k$ converging towards zero such that

$$\lim_{k \rightarrow +\infty} \left(E_1^{\varepsilon_k}, E_2^{\varepsilon_k}, \frac{B_3^{\varepsilon_k} - B_{0,3}}{\varepsilon_k} \right) = (E_1, E_2, b_3),$$

weakly in $L^2([0, T] \times \mathbb{T}^2)^3$, $\forall T > 0$. We use also the conservations of the mass and momentum

$$\partial_t \rho^\varepsilon + \operatorname{div}_x J^\varepsilon = 0, \quad (26)$$

$$\varepsilon^2 \partial_t \int_{\mathbb{R}^2} u F^\varepsilon \, du + \varepsilon^2 \operatorname{div}_x \int_{\mathbb{R}^2} (u \otimes v^\varepsilon(u)) F^\varepsilon \, du + \rho^\varepsilon(t, x) E^\varepsilon(t, x) + B_3^\varepsilon(t, x)^\perp J^\varepsilon(t, x) = 0. \quad (27)$$

By equation (25) we deduce that $\lim_{\varepsilon \searrow 0} B_3^\varepsilon(t, \cdot) = B_{0,3}$ in $L^2(\mathbb{T}^2)$ uniformly with respect to $t \in \mathbb{R}_+$ and thus from (27) we expect that at the limit for $\varepsilon \searrow 0$ one gets

$$\rho(t, x) E(t, x) + B_{0,3}^\perp J(t, x) = 0.$$

Moreover from (21) and (25) we deduce that

$$\partial_{x_1} E_2 - \partial_{x_2} E_1 = 0. \quad (28)$$

Combining with the continuity equation (26) and (22) we obtain the limit system

$$J = \rho \frac{\perp E}{B_{0,3}}, \quad \operatorname{div}_x^\perp E = 0, \quad (29)$$

$$\partial_t \rho + \operatorname{div}_x \left(\rho \frac{\perp E}{B_{0,3}} \right) = 0, \quad (30)$$

$$\operatorname{div}_x E = 1 - \rho(t, x). \quad (31)$$

The above equations can be written

$$\partial_t(\rho - 1) + \frac{{}^\perp E}{B_{0,3}} \cdot \nabla_x(\rho - 1) = 0,$$

$$\rho(t, x) - 1 = -\operatorname{div}_x E = \partial_{x_1}({}^\perp E)_2 - \partial_{x_2}({}^\perp E)_1, \quad \operatorname{div}_x {}^\perp E = 0.$$

We recognize here the Euler equations written in the so-called vorticity formulation with $\rho - 1$ standing for the vorticity and the velocity ${}^\perp E$. For the existence theory of classical solutions to the equations of ideal fluid flow we refer to [18], [23], [5], [20], [21].

The previous equations are supplemented with the initial condition E_0 given in H4, the initial condition for ρ being $\rho_0 = 1 - \operatorname{div}_x E_0$. It is easily seen by standard computation that $\frac{1}{2} \int_{\mathbb{T}^2} |E(t, x)|^2 dx$ is preserved in time. Notice also that when ε goes to zero we expect that the total kinetic and electric energy $\int_{\mathbb{T}^2} \int_{\mathbb{R}^2} ((1 + |p|^2)^{\frac{1}{2}} - 1) f^\varepsilon dp dx + \frac{1}{2} \int_{\mathbb{T}^2} |E^\varepsilon|^2 dx$ is conserved since $\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \approx 0$ (actually this happens for the Vlasov-Poisson equations with strong external magnetic field $\frac{B_{0,3}}{\varepsilon}$). Therefore we can interpret the hypotheses H2, H4 as follows: as ε goes to zero the total kinetic and electric energy of the conditions $(f_0^\varepsilon, E_0^\varepsilon)$ converges towards the electric energy of E_0 and the magnetic energy of $\frac{B_{0,3}^\varepsilon - B_{0,3}}{\varepsilon}$ goes to zero. We will see that under these hypotheses we can prove strong convergences in L^2 for the fields and also convergences in distributions sense for the charge and current densities. The same limit system has been obtained in [11], [3].

The limit system can be solved explicitly when the initial conditions depend only on x_1 . This situation arises when studying the Vlasov-Maxwell equations in the one and one-half dimensional setting [9] that is, the particle density depends on t, x_1, p_1, p_2 and the fields depend on t, x_1 . We are looking now for solutions of the limit system depending only on t and x_1 . For any 1-periodic function $u = u(x_1)$ we denote by $\langle u \rangle$ its average over one period $\langle u \rangle := \int_{\mathbb{T}^1} u(x_1) dx_1$. From (28) we deduce that $\partial_{x_1} E_2 = 0$. Integrating now (20) with respect to $x_1 \in \mathbb{T}^1$ yields

$$\frac{d}{dt} \int_{\mathbb{T}^1} E_2^\varepsilon(t, x_1) dx_1 = \int_{\mathbb{T}^1} J_2^\varepsilon(t, x_1) dx_1 = \langle J_2^\varepsilon(t) \rangle,$$

and after passing to the limit we expect that

$$\frac{d}{dt} E_2(t) = \int_{\mathbb{T}^1} J_2(t, x_1) dx_1 = -\frac{\langle \rho(t) E_1(t) \rangle}{B_{0,3}}. \quad (32)$$

Combining (19), (22) we find

$$\partial_t E_1 = \lim_{\varepsilon \searrow 0} J_1^\varepsilon(t, x_1) = \frac{\rho(t, x_1) E_2(t)}{B_{0,3}} = (1 - \partial_{x_1} E_1) \frac{E_2(t)}{B_{0,3}},$$

implying that

$$\partial_t E_1 + \frac{E_2(t)}{B_{0,3}} \partial_{x_1} E_1 = \frac{E_2(t)}{B_{0,3}}. \quad (33)$$

In this case the continuity equation becomes

$$\partial_t \rho + \frac{E_2(t)}{B_{0,3}} \partial_{x_1} \rho = 0. \quad (34)$$

Multiplying (33) by ρ and (34) by E_1 one gets

$$\partial_t(\rho E_1) + \frac{E_2(t)}{B_{0,3}} \partial_{x_1}(\rho E_1) = \frac{E_2(t)}{B_{0,3}} \rho(t, x_1),$$

and therefore by taking the average we obtain

$$\frac{d}{dt} \langle \rho(t) E_1(t) \rangle = \frac{E_2(t)}{B_{0,3}} \langle \rho(t) \rangle = \frac{E_2(t)}{B_{0,3}}. \quad (35)$$

Therefore (32), (35) can be solved with respect to E_2 and $\langle \rho E_1 \rangle$ and thus

$$E_2(t) = E_{0,2} \cos\left(\frac{t}{B_{0,3}}\right) - \langle \rho_0 E_{0,1} \rangle \sin\left(\frac{t}{B_{0,3}}\right).$$

By H3, H4 E_0 satisfies $\operatorname{div}_x^\perp E_0 = 0$, saying that indeed $E_{0,2}$ do not depend on x_1 , and $\operatorname{div}_x E_0 = 1 - \rho_0$ which becomes

$$\partial_{x_1} E_{0,1} = 1 - \rho_0(x_1), \quad x_1 \in \mathbb{T}^1.$$

Multiplying by $E_{0,1}$ and integrating over \mathbb{T}^1 yields $\langle \rho_0 E_{0,1} \rangle = \langle E_{0,1} \rangle$ and we can eliminate the function ρ_0 in the expression of E_2

$$E_2(t) = E_{0,2} \cos\left(\frac{t}{B_{0,3}}\right) - \langle E_{0,1} \rangle \sin\left(\frac{t}{B_{0,3}}\right).$$

The other unknowns can be easily expressed in terms of the characteristics $X(s; t, x_1)$ associated to $\frac{E_2}{B_{0,3}}$

$$\frac{d}{ds} X(s; t, x_1) = \frac{E_2(s)}{B_{0,3}}, \quad X(t; t, x_1) = x_1,$$

given by

$$X(s; t, x_1) = x_1 + E_{0,2} \left\{ \sin\left(\frac{s}{B_{0,3}}\right) - \sin\left(\frac{t}{B_{0,3}}\right) \right\} + \langle E_{0,1} \rangle \left\{ \cos\left(\frac{s}{B_{0,3}}\right) - \cos\left(\frac{t}{B_{0,3}}\right) \right\}.$$

Finally we obtain from (34), (33) $\rho(t, x_1) = \rho_0(X(0; t, x_1))$ and

$$E_1(t, x_1) = E_{0,1}(X(0; t, x_1)) + E_{0,2} \sin\left(\frac{t}{B_{0,3}}\right) + \langle E_{0,1} \rangle \left(\cos\left(\frac{t}{B_{0,3}}\right) - 1 \right).$$

2.2 Evolution of the modulated energy

In this paragraph we consider smooth solutions $(f^\varepsilon, E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)_{\varepsilon>0}$ for the two dimensional relativistic Vlasov-Maxwell system associated to smooth initial conditions $(F_0^\varepsilon, E_{0,1}^\varepsilon, E_{0,2}^\varepsilon, B_{0,3}^\varepsilon)_{\varepsilon>0}$ satisfying the hypotheses H1-H4. We assume also that the limit system (29), (30), (31) has a smooth solution (ρ, J, E) . Notice that the solution of the limit system satisfies $\operatorname{div}_x(\partial_t E - J) = 0$ and therefore there is a periodic function $A_3 = A_3(t, x)$ such that

$$\partial_t E_1 - \partial_{x_2} A_3 = J_1, \quad \partial_t E_2 + \partial_{x_1} A_3 = J_2. \quad (36)$$

Actually A_3 solves the elliptic space periodic problem

$$-\Delta_x A_3 = \partial_{x_2} J_1 - \partial_{x_1} J_2,$$

which has a unique periodic solution, up to an additive constant. In order to fix the constant we choose the solution with zero space average

$$\int_{\mathbb{T}^2} A_3(t, x) dx = 0, \quad t \in \mathbb{R}_+.$$

We assume that A_3 is smooth. We introduce now the modulated energy

$$\begin{aligned} \mathcal{H}^\varepsilon(t) &= \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \left(\mathcal{E}^\varepsilon(u) - D \cdot u + \frac{|D|^2}{2} \right) F^\varepsilon(t, x, u) du dx \\ &+ \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon(t, x) - E(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx, \end{aligned}$$

where $D(t, x) = \frac{J(t, x)}{\rho(t, x)} = \frac{\perp E(t, x)}{B_{0,3}}$. We intend to study the time evolution of \mathcal{H}^ε . For this we multiply the Vlasov equation (18) by the smooth function $h^\varepsilon(t, x, u) = \mathcal{E}^\varepsilon(u) - D(t, x) \cdot u + \frac{|D(t, x)|^2}{2}$. We perform our computations in several steps by observing that the Vlasov equation can be written

$$\begin{aligned} \varepsilon^2 (\partial_t F^\varepsilon + \operatorname{div}_x(v^\varepsilon(u) F^\varepsilon)) - \operatorname{div}_u \left((E(t, x) + B_{0,3} \perp v^\varepsilon(u)) F^\varepsilon \right) \\ - \operatorname{div}_u \left((E^\varepsilon(t, x) - E(t, x) + (B_3^\varepsilon(t, x) - B_{0,3}) \perp v^\varepsilon(u)) F^\varepsilon \right) = 0. \end{aligned} \quad (37)$$

Lemma 2.1 *For any $0 < \varepsilon < 1$, $T \in \mathbb{R}_+$, $t \in [0, T]$ we have the inequality*

$$\begin{aligned} (1 - \varepsilon C) \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon du dx - \varepsilon C &\leq \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} h^\varepsilon F^\varepsilon du dx \\ &\leq (1 + \varepsilon C) \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon du dx + \varepsilon C, \end{aligned} \quad (38)$$

where C is a constant depending on $\|D\|_{L^\infty([0, T] \times \mathbb{T}^2)}$. In particular

$$\varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon du dx - \varepsilon \tilde{C} \leq \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} h^\varepsilon F^\varepsilon du dx \leq \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon du dx + \varepsilon \tilde{C}.$$

Proof. We use the inequality

$$\varepsilon|u| \leq \varepsilon^2 \mathcal{E}^\varepsilon(u) + 1, \quad \forall \varepsilon > 0. \quad (39)$$

Therefore we can write

$$\begin{aligned} \varepsilon^2 \left| \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} D \cdot u F^\varepsilon \, du \, dx \right| &\leq \varepsilon \|D\|_{L^\infty} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} (\varepsilon^2 \mathcal{E}^\varepsilon(u) + 1) F^\varepsilon \, du \, dx \\ &= \varepsilon \|D\|_{L^\infty} \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon \, du \, dx + \varepsilon \|D\|_{L^\infty}, \end{aligned}$$

implying that

$$\begin{aligned} \left| \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} h^\varepsilon F^\varepsilon \, du \, dx - \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon \, du \, dx \right| &\leq \varepsilon^2 \left| \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} D \cdot u F^\varepsilon \, du \, dx \right| \\ &\quad + \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{|D|^2}{2} F^\varepsilon \, du \, dx \\ &\leq \varepsilon \|D\|_{L^\infty} \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon \, du \, dx \\ &\quad + \varepsilon \|D\|_{L^\infty} + \frac{\varepsilon^2}{2} \|D\|_{L^\infty}^2, \end{aligned}$$

and the first statement follows. The second one comes easily by using also the total energy conservation (25). \square

Lemma 2.2 For any $0 < \varepsilon < 1$, $T \in \mathbb{R}_+$, $t \in [0, T]$ we have

$$\int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \varepsilon^2 (\partial_t F^\varepsilon + \operatorname{div}_x (v^\varepsilon(u) F^\varepsilon)) h^\varepsilon \, du \, dx = \varepsilon^2 \frac{d}{dt} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} h^\varepsilon F^\varepsilon \, du \, dx - Q_1(t), \quad (40)$$

where

$$|Q_1(t)| \leq C \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon(t, x, u) \, du \, dx + C \varepsilon,$$

for some constant depending on $\|D\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$.

Proof. Integrating by parts with respect to x we deduce that the term $Q_1(t)$ in (40) has the form

$$Q_1(t) = \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} F^\varepsilon (\partial_t h^\varepsilon + v^\varepsilon(u) \cdot \nabla_x h^\varepsilon) \, du \, dx.$$

Notice that we have the inequality

$$\mathcal{E}^\varepsilon(u) \geq \frac{|u|^2}{2(1 + \varepsilon^2|u|^2)^{\frac{1}{2}}} = \frac{|u||v^\varepsilon(u)|}{2}, \quad \forall \varepsilon > 0, \quad (41)$$

and therefore we deduce

$$\begin{aligned}
 |Q_1(t)| &\leq \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} (|\partial_t D| + |(\nabla_x D)v^\varepsilon(u)|)(|u| + |D|)F^\varepsilon \, du \, dx \\
 &\leq \varepsilon^2 \|\partial_t D\|_{L^\infty} \|D\|_{L^\infty} + \varepsilon \|\partial_t D\|_{L^\infty} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} (\varepsilon^2 \mathcal{E}^\varepsilon(u) + 1)F^\varepsilon \, du \, dx \\
 &+ \varepsilon \|\nabla_x D\|_{L^\infty} \|D\|_{L^\infty} + 2\varepsilon^2 \|\nabla_x D\|_{L^\infty} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u)F^\varepsilon \, du \, dx \\
 &\leq C\varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u)F^\varepsilon \, du \, dx + C\varepsilon.
 \end{aligned}$$

□

Lemma 2.3 For any $\varepsilon > 0$, $t \in \mathbb{R}_+$ we have

$$- \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \operatorname{div}_u ((E + B_{0,3} \cdot v^\varepsilon(u))F^\varepsilon) h^\varepsilon(t, x, u) \, du \, dx = 0. \quad (42)$$

Proof. We have

$$\nabla_u h^\varepsilon = v^\varepsilon(u) - D \cdot (\cdot D \cdot v^\varepsilon(u)) = \left(-\frac{E}{B_{0,3}} \cdot v^\varepsilon(u) \right) = -\frac{1}{B_{0,3}} \cdot (E + B_{0,3} \cdot v^\varepsilon(u)),$$

and our conclusion follows easily by integration by parts. □

Lemma 2.4 For any $0 < \varepsilon < 1$, $T \in \mathbb{R}_+$, $t \in [0, T]$ we have

$$\begin{aligned}
 &- \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \operatorname{div}_u [(E^\varepsilon - E + (B_3^\varepsilon - B_{0,3}) \cdot v^\varepsilon(u))F^\varepsilon] h^\varepsilon \, du \, dx \\
 &= \frac{1}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx \\
 &- \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} [A_3 + D \cdot (E^\varepsilon - E)] \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx - Q_2(t), \quad (43)
 \end{aligned}$$

where Q_2 satisfies

$$|Q_2(t)| \leq C\varepsilon^2 + C \int_{\mathbb{T}^2} \frac{1}{2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx,$$

for some constant C depending on $\|A_3\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$, $\|D\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$.

Proof. After integration by parts with respect to u one gets

$$\begin{aligned}
& - \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \operatorname{div}_u [(E^\varepsilon - E + (B_3^\varepsilon - B_{0,3})^\perp v^\varepsilon(u)) F^\varepsilon] h^\varepsilon \, du \, dx \quad (44) \\
& = \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} F^\varepsilon (E^\varepsilon - E + (B_3^\varepsilon - B_{0,3})^\perp v^\varepsilon(u)) (v^\varepsilon(u) - D) \, du \, dx \\
& = \int_{\mathbb{T}^2} (E^\varepsilon - E) (J^\varepsilon - \rho^\varepsilon D) \, dx - \int_{\mathbb{T}^2} {}^\perp J^\varepsilon \cdot D (B_3^\varepsilon - B_{0,3}) \, dx \\
& = \int_{\mathbb{T}^2} (E^\varepsilon - E) \cdot (J^\varepsilon - J) \, dx - \int_{\mathbb{T}^2} D \cdot [(E^\varepsilon - E) (\rho^\varepsilon - \rho) + {}^\perp (J^\varepsilon - J) (B_3^\varepsilon - B_{0,3})] \, dx,
\end{aligned}$$

since $D \cdot {}^\perp J = 0$. Combining (19), (20), (21), (36) yields

$$\partial_t (E_1^\varepsilon - E_1) - \frac{1}{\varepsilon} \partial_{x_2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) + \partial_{x_2} A_3 = J_1^\varepsilon - J_1, \quad (45)$$

$$\partial_t (E_2^\varepsilon - E_2) + \frac{1}{\varepsilon} \partial_{x_1} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) - \partial_{x_1} A_3 = J_2^\varepsilon - J_2, \quad (46)$$

$$\partial_t \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) + \frac{1}{\varepsilon} \partial_{x_1} (E_2^\varepsilon - E_2) - \frac{1}{\varepsilon} \partial_{x_2} (E_1^\varepsilon - E_1) = 0. \quad (47)$$

Notice that in the last equation we have used $\partial_{x_1} E_2 - \partial_{x_2} E_1 = 0$. Multiplying (45) by $E_1^\varepsilon - E_1$, (46) by $E_2^\varepsilon - E_2$ and (47) by $\varepsilon^{-1} (B_3^\varepsilon - B_{0,3})$ implies

$$\begin{aligned}
\frac{1}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} \, dx + \int_{\mathbb{T}^2} \{ (E_1^\varepsilon - E_1) \partial_{x_2} A_3 - (E_2^\varepsilon - E_2) \partial_{x_1} A_3 \} \, dx \\
= \int_{\mathbb{T}^2} (J^\varepsilon - J) \cdot (E^\varepsilon - E) \, dx.
\end{aligned}$$

Using one more time (21) we can write

$$\begin{aligned}
\int_{\mathbb{T}^2} (E_1^\varepsilon \partial_{x_2} A_3 - E_2^\varepsilon \partial_{x_1} A_3) \, dx &= \int_{\mathbb{T}^2} A_3 (\partial_{x_1} E_2^\varepsilon - \partial_{x_2} E_1^\varepsilon) \, dx \\
&= - \int_{\mathbb{T}^2} A_3 \partial_t (B_3^\varepsilon - B_{0,3}) \, dx \\
&= -\varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx + \varepsilon \int_{\mathbb{T}^2} \partial_t A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx.
\end{aligned}$$

Since $\int_{\mathbb{T}^2} (E_1 \partial_{x_2} A_3 - E_2 \partial_{x_1} A_3) \, dx = 0$ finally one gets

$$\begin{aligned}
\int_{\mathbb{T}^2} (E^\varepsilon - E) \cdot (J^\varepsilon - J) \, dx &= \frac{1}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} \, dx \quad (48) \\
&- \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx + \varepsilon \int_{\mathbb{T}^2} \partial_t A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx.
\end{aligned}$$

We transform now the last term in (44) using (45), (46), (47) and (22), (31). We have

$$\begin{aligned}
 & - (\rho^\varepsilon - \rho)(E^\varepsilon - E) - (B_3^\varepsilon - B_{0,3})^\perp (J^\varepsilon - J) = \operatorname{div}_x (E^\varepsilon - E) (E^\varepsilon - E) \\
 & - (B_3^\varepsilon - B_{0,3}) \left\{ \partial_t^\perp (E^\varepsilon - E) + \frac{1}{\varepsilon} \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) - \nabla_x A_3 \right\} \\
 & = \operatorname{div}_x (E^\varepsilon - E) (E^\varepsilon - E) + \partial_t (B_3^\varepsilon - B_{0,3})^\perp (E^\varepsilon - E) - \partial_t ((B_3^\varepsilon - B_{0,3})^\perp (E^\varepsilon - E)) \\
 & - \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) + (B_3^\varepsilon - B_{0,3}) \nabla_x A_3 \\
 & = \operatorname{div}_x (E^\varepsilon - E) (E^\varepsilon - E) - \operatorname{div}_x^\perp (E^\varepsilon - E)^\perp (E^\varepsilon - E) - \frac{1}{2} \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \\
 & - \varepsilon \partial_t \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon}^\perp (E^\varepsilon - E) \right) + \varepsilon \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \nabla_x A_3. \tag{49}
 \end{aligned}$$

We have the identity

$$\operatorname{div}_x w w - \operatorname{div}_x^\perp w^\perp w = \operatorname{div}_x (w \otimes w) - \frac{1}{2} \nabla_x |w|^2,$$

for any $w = (w_1, w_2) \in C^1(\mathbb{R}^2)^2$. Multiplying (49) by D and integrating by parts with respect to x yields

$$\begin{aligned}
 & - \int_{\mathbb{T}^2} D \cdot [(\rho^\varepsilon - \rho)(E^\varepsilon - E) + (B_3^\varepsilon - B_{0,3})^\perp (J^\varepsilon - J)] dx \\
 & = - \int_{\mathbb{T}^2} (\nabla_x D (E^\varepsilon - E)) (E^\varepsilon - E) dx + \frac{1}{2} \int_{\mathbb{T}^2} (\operatorname{div}_x D) \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx \\
 & - \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} D \cdot^\perp (E^\varepsilon - E) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx + \varepsilon \int_{\mathbb{T}^2} \partial_t D \cdot^\perp (E^\varepsilon - E) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx \\
 & + \varepsilon \int_{\mathbb{T}^2} (D \cdot \nabla_x A_3) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx. \tag{50}
 \end{aligned}$$

Combining (44), (48), (50) and observing that $\operatorname{div}_x D = 0$ we deduce that the term $Q_2(t)$ in (43) has the form

$$\begin{aligned}
 -Q_2(t) & = \varepsilon \int_{\mathbb{T}^2} (\partial_t A_3 + D \cdot \nabla_x A_3) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx + \varepsilon \int_{\mathbb{T}^2} \partial_t D \cdot^\perp (E^\varepsilon - E) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx \\
 & - \int_{\mathbb{T}^2} (\nabla_x D (E^\varepsilon - E)) \cdot (E^\varepsilon - E) dx.
 \end{aligned}$$

By using the inequality

$$\varepsilon \left| (\partial_t A_3 + D \cdot \nabla_x A_3) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \right| \leq \frac{\varepsilon^2}{2} |\partial_t A_3 + D \cdot \nabla_x A_3|^2 + \frac{1}{2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2,$$

we obtain

$$|Q_2(t)| \leq C\varepsilon^2 + C \int_{\mathbb{T}^2} \frac{1}{2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx,$$

for some constant depending on $\|A_3\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$, $\|D\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$. \square

Proposition 2.1 *There is a constant C depending on $\|D\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$, $\|A_3\|_{W^{1,\infty}([0,T] \times \mathbb{T}^2)}$ such that for any $0 < \varepsilon < \varepsilon(C)$, $t \in [0, T]$ we have*

$$\begin{aligned} & \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon(t, x, u) du dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon(t, x) - E(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx \\ & \leq C_1(t) \left\{ \varepsilon + \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F_0^\varepsilon du dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E_0^\varepsilon - E_0|^2 + \left(\frac{B_{0,3}^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx \right\}, \end{aligned}$$

where $C_1(t) = (3 + 2C(4+t))e^{2Ct}$, $t \in [0, T]$.

Proof. Using the Lemmas 2.2, 2.3, 2.4 in (37) yields

$$\frac{d}{dt} \mathcal{H}^\varepsilon(t) - \varepsilon \frac{d}{dt} \mathcal{R}^\varepsilon(t) = Q_1(t) + Q_2(t) \leq C\varepsilon + C\mathcal{W}^\varepsilon(t), \quad (51)$$

where

$$\mathcal{R}^\varepsilon(t) = \int_{\mathbb{T}^2} (A_3 + D \cdot^\perp (E^\varepsilon - E)) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx,$$

and

$$\mathcal{W}^\varepsilon(t) = \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \mathcal{E}^\varepsilon(u) F^\varepsilon du dx + \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right\} dx.$$

By Lemma 2.1 we deduce that $|\mathcal{H}^\varepsilon(t) - \mathcal{W}^\varepsilon(t)| \leq C\varepsilon$. Observe also that we have $|\mathcal{R}^\varepsilon(t)| \leq C(1 + \mathcal{W}^\varepsilon(t))$. Integrating (51) over $[0, t]$ one gets

$$\mathcal{H}^\varepsilon(t) - \varepsilon \mathcal{R}^\varepsilon(t) \leq \mathcal{H}^\varepsilon(0) - \varepsilon \mathcal{R}^\varepsilon(0) + C\varepsilon t + C \int_0^t \mathcal{W}^\varepsilon(s) ds, \quad t \in [0, T]. \quad (52)$$

Notice that for any $\varepsilon < 1/(2C)$ we have

$$\mathcal{H}^\varepsilon(t) - \varepsilon \mathcal{R}^\varepsilon(t) \geq \mathcal{W}^\varepsilon(t) - 2C\varepsilon - \varepsilon C\mathcal{W}^\varepsilon(t) \geq \frac{1}{2} \mathcal{W}^\varepsilon(t) - 2C\varepsilon,$$

and

$$\mathcal{H}^\varepsilon(0) - \varepsilon \mathcal{R}^\varepsilon(0) \leq \mathcal{W}^\varepsilon(0) + 2C\varepsilon + \varepsilon C\mathcal{W}^\varepsilon(0) \leq \frac{3}{2} \mathcal{W}^\varepsilon(0) + 2C\varepsilon.$$

Combining the above inequalities with (52) implies

$$\frac{1}{2} \mathcal{W}^\varepsilon(t) \leq \frac{3}{2} \mathcal{W}^\varepsilon(0) + C\varepsilon(4+t) + C \int_0^t \mathcal{W}^\varepsilon(s) ds,$$

and the conclusion follows easily by Gronwall lemma. \square

The estimate for the modulated energy allows us to justify the convergence of the Vlasov-Maxwell system with strong initial magnetic field towards the vorticity formulation of the Euler equations.

Theorem 2.1 *Assume that the initial conditions $(f_0^\varepsilon, E_{0,1}^\varepsilon, E_{0,2}^\varepsilon, B_{0,3}^\varepsilon)_{\varepsilon>0}$ are smooth and satisfy the hypotheses H1-H4. We denote by $(f^\varepsilon, E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)_{\varepsilon>0}$ the solutions of the two dimensional problems (9), (10), (11), (12), (13), (14) and we suppose that the limit system (29), (30), (31) corresponding to the initial electric field E_0 appearing in H4 has a smooth solution (ρ, J, E) . Then for any $T \in \mathbb{R}_+$ we have*

$$\begin{aligned} \lim_{\varepsilon \searrow 0} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} ((1 + |p|^2)^{\frac{1}{2}} - 1) f^\varepsilon(t, x, p) dp dx + \frac{1}{2} \int_{\mathbb{T}^2} |E^\varepsilon(t, x) - E(t, x)|^2 dx \\ + \frac{1}{2} \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 dx = 0, \quad \text{uniformly for } t \in [0, T], \\ \lim_{\varepsilon \searrow 0} \rho^\varepsilon = \rho, \quad \lim_{\varepsilon \searrow 0} \frac{j^\varepsilon}{\varepsilon} = J \quad \text{in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2). \end{aligned}$$

Proof. The first statement comes by Proposition 2.1. The convergence of the charge densities $(\rho^\varepsilon)_{\varepsilon>0}$ follows easily by (22), (31) since for any $\varphi \in C_c^1(\mathbb{R}_+ \times \mathbb{T}^2)$ we have

$$\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (\rho^\varepsilon - \rho)(t, x) \varphi(t, x) dx dt = \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (E^\varepsilon - E)(t, x) \cdot \nabla_x \varphi dx dt = 0.$$

For the convergence of the current densities $(J^\varepsilon)_{\varepsilon>0} = \left(\frac{j^\varepsilon}{\varepsilon} \right)_{\varepsilon>0}$ we use the momentum conservation (27). It is easily seen by Proposition 2.1 and the inequalities (39), (41) that

$$\begin{aligned} \lim_{\varepsilon \searrow 0} \varepsilon^2 \partial_t \int_{\mathbb{R}^2} u F^\varepsilon du = 0 \quad \text{in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2, \\ \lim_{\varepsilon \searrow 0} \varepsilon^2 \operatorname{div}_x \int_{\mathbb{R}^2} (u \otimes v^\varepsilon(u)) F^\varepsilon du = 0 \quad \text{in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2. \end{aligned}$$

We introduce the quadratic form $\mathcal{F}(w) = \operatorname{div}_x w w - \operatorname{div}_x^\perp w^\perp w$. By using (19), (20), (21), (22) one gets

$$\operatorname{div}_x E^\varepsilon E^\varepsilon - (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon = \mathcal{F}(E^\varepsilon) - \frac{1}{2} \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 - \varepsilon \partial_t \left\{ \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^\perp E^\varepsilon \right\}.$$

Notice that $\mathcal{F}(E^\varepsilon) = \operatorname{div}_x (E^\varepsilon \otimes E^\varepsilon) - \frac{1}{2} \nabla_x |E^\varepsilon|^2$ and since we know that $(E^\varepsilon)_{\varepsilon>0}$ converges strongly towards E in $L^2([0, T] \times \mathbb{T}^2)^2$ we deduce that

$$\lim_{\varepsilon \searrow 0} \mathcal{F}(E^\varepsilon) = \mathcal{F}(E) = \operatorname{div}_x E E \quad \text{in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2.$$

Similarly, as $\lim_{\varepsilon \searrow 0} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) = 0$ in $L^2([0, T] \times \mathbb{T}^2)$ we have

$$\lim_{\varepsilon \searrow 0} \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 = \lim_{\varepsilon \searrow 0} \varepsilon \partial_t \left\{ \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^\perp E^\varepsilon \right\} = 0 \text{ in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2),$$

and therefore

$$\lim_{\varepsilon \searrow 0} \{ \operatorname{div}_x E^\varepsilon E^\varepsilon - (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon \} = \operatorname{div}_x E E \text{ in } \mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2.$$

Passing now to the limit in (27) one gets in $\mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2$

$$\begin{aligned} \lim_{\varepsilon \searrow 0} B_{0,3}^\perp J^\varepsilon &= - \lim_{\varepsilon \searrow 0} \{ \rho^\varepsilon E^\varepsilon + (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon \} \\ &= - \lim_{\varepsilon \searrow 0} E^\varepsilon + \lim_{\varepsilon \searrow 0} \{ \operatorname{div}_x E^\varepsilon E^\varepsilon - (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon \} \\ &= -E + \operatorname{div}_x E E = -\rho E, \end{aligned}$$

saying that $\lim_{\varepsilon \searrow 0} J^\varepsilon = \rho \frac{\perp E}{B_{0,3}} = J$ in $\mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)^2$. \square

We end this section with a convergence result in distribution sense for $\left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)_{\varepsilon > 0}$.

Corollary 2.1 *Besides the hypotheses of Theorem 2.1 assume that the following condition holds*

$$H5) \lim_{\varepsilon \searrow 0} \int_{\mathbb{T}^2} \frac{B_{0,3}^\varepsilon(x) - B_{0,3}}{\varepsilon^2} dx = 0.$$

Then we have the convergence $\lim_{\varepsilon \searrow 0} \frac{B_3^\varepsilon - B_{0,3}}{\varepsilon^2} = A_3$ in $\mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)$.

Proof. Combining (45), (46) we have

$$\partial_t (E^\varepsilon - E) - \perp \nabla_x (A_3^\varepsilon - A_3) = J^\varepsilon - J,$$

where $A_3^\varepsilon := \frac{B_3^\varepsilon - B_{0,3}}{\varepsilon^2}$. We deduce easily that

$$\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \nabla_x \varphi dx dt = 0, \forall \varphi \in C_c^1(\mathbb{R}_+ \times \mathbb{T}^2).$$

In particular we have $\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \operatorname{div}_x \varphi dx dt = 0$ for any $\varphi \in C_c^1(\mathbb{R}_+ \times \mathbb{T}^2)^2$.

Take now $\psi \in C_c^\infty(\mathbb{R}_+ \times \mathbb{T}^2)$ satisfying $\int_{\mathbb{T}^2} \psi(t, x) dx = 0$, $t \in \mathbb{R}_+$ and denote by u the solution of $-\Delta_x u(t) = \psi(t, x)$, $x \in \mathbb{T}^2$, $t \in \mathbb{R}_+$, verifying $\int_{\mathbb{T}^2} u(t, x) dx = 0$, $t \in \mathbb{R}_+$. We have

$$\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \psi(t, x) dx dt = - \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \operatorname{div}_x (\nabla_x u) dx dt = 0.$$

Take now $\psi \in C_c^\infty(\mathbb{R}_+ \times \mathbb{T}^2)$ and observing that $\psi - \int_{\mathbb{T}^2} \psi dx$ has zero space average we obtain

$$\begin{aligned} \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \psi(t, x) dx dt &= \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) (\psi - \langle \psi \rangle + \langle \psi \rangle) dx dt \\ &= \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \langle \psi(t) \rangle \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) dx dt. \end{aligned}$$

Recall that by definition $\int_{\mathbb{T}^2} A_3(t, x) dx = 0$, $t \in \mathbb{R}_+$ and by integrating (47) we deduce that $\frac{d}{dt} \int_{\mathbb{T}^2} A_3^\varepsilon(t, x) dx = 0$. Therefore the hypothesis H5 yields

$$\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \int_{\mathbb{T}^2} (A_3^\varepsilon - A_3) \psi(t, x) dx dt = \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}_+} \langle \psi(t) \rangle \int_{\mathbb{T}^2} A_3^\varepsilon(0, x) dx dt = 0.$$

□

3 Other systems

We can consider the Vlasov-Maxwell system in the non relativistic setting. In the two dimensional case the Vlasov equation (9) becomes

$$\partial_t f^\varepsilon + \frac{p}{\varepsilon} \cdot \nabla_x f^\varepsilon - \frac{1}{\varepsilon} \left(E^\varepsilon(t, x) + \frac{B_3^\varepsilon(t, x)}{\varepsilon} \perp p \right) \cdot \nabla_p f^\varepsilon = 0, \quad (t, x, p) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2, \quad (53)$$

where $(E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)_{\varepsilon > 0}$ solve the two dimensional Maxwell equations with the charge density $\int_{\mathbb{R}^2} f^\varepsilon dp$ and the current density $\int_{\mathbb{R}^2} p f^\varepsilon dp$. Rescaling the momentum by $p = \varepsilon u$ and the particle density by $F^\varepsilon(t, x, u) = \varepsilon^2 f^\varepsilon(t, x, \varepsilon u)$ leads to the same equations as those in (18), (19), (20), (21), (22) with $v^\varepsilon(u)$ replaced by u and $J^\varepsilon(t, x)$ replaced by $\int_{\mathbb{R}^2} u F^\varepsilon(t, x, u) du$. We assume that the hypotheses H1, H3, H4 hold and we replace H2 by

$$\lim_{\varepsilon \searrow 0} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{|p|^2}{2} f_0^\varepsilon(x, p) dp dx = 0,$$

or equivalently by

$$\lim_{\varepsilon \searrow 0} \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{|u|^2}{2} F_0^\varepsilon(x, u) du dx = 0.$$

Following the previous method we show the convergence towards a solution (ρ, J, E) of (29), (30), (31). The modulated energy is given by

$$\begin{aligned} \mathcal{H}_2^\varepsilon(t) &= \varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{1}{2} |u - D(t, x)|^2 F^\varepsilon(t, x, u) du dx \\ &+ \frac{1}{2} \int_{\mathbb{T}^2} \left\{ |E^\varepsilon(t, x) - E(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right\} dx, \end{aligned}$$

where $D(t, x) = \frac{J(t, x)}{\rho(t, x)} = \frac{\perp E(t, x)}{B_{0,3}}$.

Proposition 3.1 *There is a constant C such that for any $\varepsilon > 0$ small enough and $t \in [0, T]$ we have*

$$\mathcal{H}_2^\varepsilon(t) \leq C(\varepsilon^2 + \mathcal{H}_2^\varepsilon(0)). \quad (54)$$

Moreover if

$$\sup_{\varepsilon>0} \varepsilon^{-2} \left\{ \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{|p|^2}{2} f_0^\varepsilon(x, p) dp dx + \frac{1}{2} \int_{\mathbb{T}^2} |E_0^\varepsilon - E_0|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 dx \right\} < +\infty,$$

then

$$\sup_{\varepsilon>0, t \in [0, T]} \varepsilon^{-2} \mathcal{H}_2^\varepsilon(t) < +\infty, \quad \forall T \in \mathbb{R}_+. \quad (55)$$

In particular

$$\sup_{\varepsilon>0} \varepsilon^{-1} \|E^\varepsilon - E\|_{L^\infty([0, T]; L^2(\mathbb{T}^2)^2)} + \sup_{\varepsilon>0} \varepsilon^{-1} \left\| \frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right\|_{L^\infty([0, T]; L^2(\mathbb{T}^2))} < +\infty,$$

and

$$\sup_{\varepsilon>0} \varepsilon^{-1} \|\rho^\varepsilon - \rho\|_{L^\infty([0, T]; H^{-1}(\mathbb{T}^2))} + \sup_{\varepsilon>0} \varepsilon^{-1} \|J^\varepsilon - J\|_{W^{-1,1}([0, T] \times \mathbb{T}^2)^2} < +\infty.$$

Proof. Let us give some details. As in Lemma 2.2, by using the inequality $|u - D| \leq 1/2 + |u - D|^2/2$ we have

$$\int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{\varepsilon^2}{2} (\partial_t F^\varepsilon + \operatorname{div}_x(u F^\varepsilon)) |u - D|^2 du dx = \frac{\varepsilon^2}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} F^\varepsilon |u - D|^2 du dx - \tilde{Q}_1(t),$$

where

$$|\tilde{Q}_1(t)| \leq C\varepsilon^2 \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \frac{1}{2} |u - D|^2 F^\varepsilon du dx + C\varepsilon^2,$$

for some constant depending on $\|D\|_{W^{1,\infty}([0, T] \times \mathbb{T}^2)^2}$. Exactly as in the relativistic case (see Lemmas 2.3, 2.4) we have

$$- \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \operatorname{div}_u ((E + B_{0,3}^\perp u) F^\varepsilon) \frac{1}{2} |u - D|^2 du dx = 0,$$

and

$$\begin{aligned} & - \int_{\mathbb{T}^2} \int_{\mathbb{R}^2} \operatorname{div}_u ((E^\varepsilon - E + (B_3^\varepsilon - B_{0,3})^\perp u) F^\varepsilon) \frac{1}{2} |u - D|^2 du dx \\ &= \frac{1}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \left(|E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right) dx \\ & - \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} (A_3 + D \cdot^\perp (E^\varepsilon - E)) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx - \tilde{Q}_2(t), \end{aligned}$$

where

$$|\tilde{Q}_2(t)| \leq C \int_{\mathbb{T}^2} \frac{1}{2} \left(|E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right) dx + C\varepsilon^2,$$

for some constant C depending on $\|A_3\|_{W^{1,\infty}([0,T[\times\mathbb{T}^2)}$, $\|D\|_{W^{1,\infty}([0,T[\times\mathbb{T}^2)^2}$. Combining the above computations yields

$$\frac{d}{dt}\mathcal{H}_2^\varepsilon(t) - \varepsilon \frac{d}{dt}\mathcal{R}_2^\varepsilon \leq C\varepsilon^2 + C\mathcal{H}_2^\varepsilon(t),$$

where $\mathcal{R}_2^\varepsilon(t) = \int_{\mathbb{T}^2} (A_3 + D \cdot^\perp (E^\varepsilon - E)) \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) dx$. The inequality (54) follows immediately by Gronwall lemma, using that

$$\begin{aligned} \varepsilon |\mathcal{R}_2^\varepsilon(t)| &\leq \varepsilon^2 \int_{\mathbb{T}^2} |A_3(t, x)|^2 dx + \frac{1}{4} \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 dx + C\varepsilon \mathcal{H}_2^\varepsilon(t) \\ &\leq \left(\frac{1}{2} + C\varepsilon \right) \mathcal{H}_2^\varepsilon(t) + C\varepsilon^2 < \frac{3}{4} \mathcal{H}_2^\varepsilon(t) + C\varepsilon^2, \end{aligned}$$

for ε small enough. The bound of $\left(\frac{\rho^\varepsilon - \rho}{\varepsilon} \right)_{\varepsilon > 0}$ in $L^\infty([0, T[; H^{-1}(\mathbb{T}^2))$ is obvious. The estimate for $\left(\frac{J^\varepsilon - J}{\varepsilon} \right)_{\varepsilon > 0}$ follows by combining the arguments in Theorem 2.1 and (55). Indeed, by the non relativistic version of (27) we have

$$\begin{aligned} B_{0,3}^\perp (J^\varepsilon - J) &= -\varepsilon^2 \left(\partial_t \int_{\mathbb{R}^2} u F^\varepsilon du + \operatorname{div}_x \int_{\mathbb{R}^2} (u \otimes u) F^\varepsilon du \right) \\ &\quad - (E^\varepsilon - E) + \operatorname{div}_x E^\varepsilon E^\varepsilon - (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon - \operatorname{div}_x E E. \end{aligned}$$

For any $\varphi \in W^{1,\infty}([0, T[\times\mathbb{T}^2)^2$ we have by (55)

$$\sup_{\varepsilon > 0} \varepsilon^2 \left| \left\langle \partial_t \int_{\mathbb{R}^2} u F^\varepsilon du + \operatorname{div}_x \int_{\mathbb{R}^2} (u \otimes u) F^\varepsilon du, \varphi \right\rangle \right| \leq C\varepsilon^2 \|\varphi\|_{W^{1,\infty}(\mathbb{R}_+ \times \mathbb{T}^2)^2},$$

and $|\langle E^\varepsilon - E, \varphi \rangle| \leq C\varepsilon \|\varphi\|_{L^\infty(\mathbb{R}_+ \times \mathbb{T}^2)^2}$. As in the proof of Theorem 2.1 we can write

$$\begin{aligned} S^\varepsilon &:= \operatorname{div}_x E^\varepsilon E^\varepsilon - (B_3^\varepsilon - B_{0,3})^\perp J^\varepsilon - \operatorname{div}_x E E \\ &= \mathcal{F}(E^\varepsilon) - \mathcal{F}(E) - \frac{1}{2} \nabla_x \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 - \varepsilon \partial_t \left\{ \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^\perp E^\varepsilon \right\}. \end{aligned}$$

It is easily seen that $|\langle S^\varepsilon, \varphi \rangle| \leq C\varepsilon \|\varphi\|_{W^{1,\infty}(\mathbb{R}_+ \times \mathbb{T}^2)^2}$. Finally one gets

$$|B_{0,3} \langle^\perp (J^\varepsilon - J), \varphi \rangle| \leq C\varepsilon \|\varphi\|_{W^{1,\infty}([0,T[\times\mathbb{T}^2)^2},$$

saying that $\sup_{\varepsilon > 0} \varepsilon^{-1} \|J^\varepsilon - J\|_{W^{-1,1}([0,T[\times\mathbb{T}^2)^2} < +\infty$. \square

Remark 3.1 *The previous result says that the solution of the limit system is a first order approximation for the non relativistic system (53), (19), (20), (21), (22), i.e.,*

$$E^\varepsilon = E + \varepsilon \mathcal{O}(\varepsilon), \quad \frac{B_3^\varepsilon}{\varepsilon} = \frac{B_{0,3}}{\varepsilon} + \varepsilon \mathcal{O}(\varepsilon), \quad \rho^\varepsilon = \rho + \varepsilon \mathcal{O}(\varepsilon), \quad J^\varepsilon = J + \varepsilon \mathcal{O}(\varepsilon),$$

in the corresponding spaces.

As in [16] we analyze also the case of distribution functions of the form

$$F^\varepsilon(t, x, u) = \rho^\varepsilon(t, x) \delta(u - u^\varepsilon(t, x)), \quad (t, x, u) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2,$$

with a macroscopic density $\rho^\varepsilon(t, x)$ and a bulk velocity $u^\varepsilon(t, x)$, or equivalently

$$f^\varepsilon(t, x, p) = \rho^\varepsilon(t, x) \delta(p - \varepsilon u^\varepsilon(t, x)), \quad (t, x, p) \in \mathbb{R}_+ \times \mathbb{T}^2 \times \mathbb{R}^2.$$

Following [6] the mass and momentum conservations lead to the equations

$$\partial_t \rho^\varepsilon + \operatorname{div}_x(\rho^\varepsilon u^\varepsilon) = 0, \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (56)$$

$$\partial_t(\rho^\varepsilon u^\varepsilon) + \operatorname{div}_x(\rho^\varepsilon(u^\varepsilon \otimes u^\varepsilon)) + \frac{1}{\varepsilon^2} \rho^\varepsilon(E^\varepsilon(t, x) + B_3^\varepsilon(t, x) \perp u^\varepsilon(t, x)) = 0, \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2, \quad (57)$$

coupled to the Maxwell equations (19), (20), (21), (22) with $J^\varepsilon(t, x) = \rho^\varepsilon(t, x) u^\varepsilon(t, x)$. By standard computations we obtain the conservation of the total energy

$$\frac{d}{dt} \left\{ \int_{\mathbb{T}^2} \frac{\varepsilon^2}{2} |u^\varepsilon(t, x)|^2 \rho^\varepsilon(t, x) dx + \frac{1}{2} \int_{\mathbb{T}^2} \left(|E^\varepsilon(t, x)|^2 + \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 \right) dx \right\} = 0. \quad (58)$$

We obtain the same limit system

$$u(t, x) = \frac{\perp E(t, x)}{B_{0,3}}, \quad \operatorname{div}_x \perp E = 0, \quad \partial_t \rho + \frac{\perp E}{B_{0,3}} \cdot \nabla_x \rho = 0, \quad \operatorname{div}_x E = 1 - \rho, \quad (t, x) \in \mathbb{R}_+ \times \mathbb{T}^2.$$

We work with smooth solutions $(\rho^\varepsilon, u^\varepsilon, E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)_{\varepsilon > 0}$, (ρ, u, E) and we define the modulated energy

$$\begin{aligned} \mathcal{H}_3^\varepsilon(t) &= \int_{\mathbb{T}^2} \frac{\varepsilon^2}{2} |u^\varepsilon(t, x) - u(t, x)|^2 \rho^\varepsilon(t, x) dx + \frac{1}{2} \int_{\mathbb{T}^2} |E^\varepsilon(t, x) - E(t, x)|^2 dx \\ &+ \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon(t, x) - B_{0,3}}{\varepsilon} \right)^2 dx. \end{aligned}$$

We study the time evolution of $\mathcal{H}_3^\varepsilon$ by using the equations for $(\rho^\varepsilon, u^\varepsilon, E_1^\varepsilon, E_2^\varepsilon, B_3^\varepsilon)$ and (ρ, u, E) . By using (56) notice that (57) can be written

$$\partial_t u^\varepsilon + (u^\varepsilon \cdot \nabla_x) u^\varepsilon + \frac{1}{\varepsilon^2} (E^\varepsilon(t, x) + B_3^\varepsilon(t, x) \perp u^\varepsilon(t, x)) = 0. \quad (59)$$

We deduce that

$$\partial_t(u^\varepsilon - u) + (u^\varepsilon \cdot \nabla_x)(u^\varepsilon - u) + \frac{1}{\varepsilon^2} (E^\varepsilon(t, x) + B_3^\varepsilon(t, x) \perp u^\varepsilon(t, x)) = -\partial_t u - (u^\varepsilon \cdot \nabla_x) u.$$

Multiplying by $\rho^\varepsilon(u^\varepsilon - u)$ yields

$$\begin{aligned} \frac{\rho^\varepsilon}{2} \partial_t |u^\varepsilon - u|^2 &+ \frac{\rho^\varepsilon}{2} (u^\varepsilon \cdot \nabla_x) |u^\varepsilon - u|^2 + \frac{\rho^\varepsilon}{\varepsilon^2} (E^\varepsilon + B_3^\varepsilon \perp u^\varepsilon) \cdot (u^\varepsilon - u) \\ &= -\rho^\varepsilon (\partial_t u + (u^\varepsilon \cdot \nabla_x) u) \cdot (u^\varepsilon - u). \end{aligned} \quad (60)$$

Adding to the above equation the equation (56) multiplied by $|u^\varepsilon - u|^2/2$ we deduce that

$$\begin{aligned} \frac{1}{2} \partial_t (\rho^\varepsilon |u^\varepsilon - u|^2) &+ \frac{1}{2} \operatorname{div}_x (\rho^\varepsilon |u^\varepsilon - u|^2 u^\varepsilon) + \frac{\rho^\varepsilon}{\varepsilon^2} (E^\varepsilon + B_3^\varepsilon \perp u^\varepsilon) \cdot (u^\varepsilon - u) \\ &= -\rho^\varepsilon (\partial_t u + (u^\varepsilon \cdot \nabla_x) u) \cdot (u^\varepsilon - u). \end{aligned} \quad (61)$$

Notice that $\perp(u^\varepsilon - u) = \perp u^\varepsilon + \frac{E}{B_{0,3}}$ and thus $(E + B_{0,3} \perp u^\varepsilon) \cdot (u^\varepsilon - u) = 0$, implying that

$$\begin{aligned} \rho^\varepsilon (E^\varepsilon + B_3^\varepsilon \perp u^\varepsilon) \cdot (u^\varepsilon - u) &= (E^\varepsilon - E) \cdot (\rho^\varepsilon u^\varepsilon - \rho u) + \operatorname{div}_x (E^\varepsilon - E) (E^\varepsilon - E) \cdot u \\ &- (B_3^\varepsilon - B_{0,3}) \perp (\rho^\varepsilon u^\varepsilon - \rho u) \cdot u. \end{aligned} \quad (62)$$

Using now the equations

$$\begin{aligned} \partial_t (E_1^\varepsilon - E_1) - \frac{1}{\varepsilon} \partial_{x_2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) + \partial_{x_2} A_3 &= \rho^\varepsilon u_1^\varepsilon - \rho u_1, \\ \partial_t (E_2^\varepsilon - E_2) + \frac{1}{\varepsilon} \partial_{x_1} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) - \partial_{x_1} A_3 &= \rho^\varepsilon u_2^\varepsilon - \rho u_2, \\ \partial_t \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) + \frac{1}{\varepsilon} \partial_{x_1} (E_2^\varepsilon - E_2) - \frac{1}{\varepsilon} \partial_{x_2} (E_1^\varepsilon - E_1) &= 0, \end{aligned}$$

one gets as before

$$\begin{aligned} \int_{\mathbb{T}^2} (E^\varepsilon - E) \cdot (\rho^\varepsilon u^\varepsilon - \rho u) \, dx &= \frac{1}{2} \frac{d}{dt} \int_{\mathbb{T}^2} \left(|E^\varepsilon - E|^2 + \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right)^2 \right) \, dx \\ - \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx &+ \varepsilon \int_{\mathbb{T}^2} \partial_t A_3 \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \, dx, \end{aligned} \quad (63)$$

and

$$\begin{aligned} &\int_{\mathbb{T}^2} \{ \operatorname{div}_x (E^\varepsilon - E) (E^\varepsilon - E) - (B_3^\varepsilon - B_{0,3}) \perp (\rho^\varepsilon u^\varepsilon - \rho u) \} \cdot u \, dx \\ &= - \int_{\mathbb{T}^2} ((\nabla_x u) (E^\varepsilon - E)) \cdot (E^\varepsilon - E) \, dx - \varepsilon \frac{d}{dt} \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \perp (E^\varepsilon - E) \cdot u \, dx \\ &+ \varepsilon \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \perp (E^\varepsilon - E) \cdot \partial_t u \, dx + \varepsilon \int_{\mathbb{T}^2} \left(\frac{B_3^\varepsilon - B_{0,3}}{\varepsilon} \right) \nabla_x A_3 \cdot u \, dx. \end{aligned} \quad (64)$$

Combining (61), (62), (63), (64) and the energy conservation (58) we obtain

$$\mathcal{H}_3^\varepsilon(t) \leq C(\varepsilon^2 + \mathcal{H}_3^\varepsilon(0)) + C \int_0^t \mathcal{H}_3^\varepsilon(s) ds,$$

implying by Gronwall lemma that $\lim_{\varepsilon \searrow 0} \mathcal{H}_3^\varepsilon(t) = 0$ uniformly on compact subsets of \mathbb{R}_+ , provided that the initial conditions satisfy the hypotheses

$$\rho_0^\varepsilon \geq 0, \quad \int_{\mathbb{T}^2} \rho_0^\varepsilon(x) dx = 1, \quad \lim_{\varepsilon \searrow 0} \varepsilon^2 \int_{\mathbb{T}^2} \frac{|u_0^\varepsilon(x)|^2}{2} \rho_0^\varepsilon(x) dx = 0,$$

and H3, H4. Therefore we deduce the convergences

$$\lim_{\varepsilon \searrow 0} \left(E_1^\varepsilon(t) - E_1(t), E_2^\varepsilon(t) - E_2(t), \frac{B_3^\varepsilon(t) - B_{0,3}}{\varepsilon} \right) = (0, 0, 0) \text{ strongly in } L^2(\mathbb{T}^2)^3,$$

uniformly for t in compact subsets of \mathbb{R}_+ . We can show as before the convergence of the charge and current densities in $\mathcal{D}'(\mathbb{R}_+ \times \mathbb{T}^2)$

$$\lim_{\varepsilon \searrow 0} \rho^\varepsilon = \rho, \quad \lim_{\varepsilon \searrow 0} (\rho^\varepsilon u^\varepsilon) = \rho u = \rho \frac{\perp E}{B_{0,3}}.$$

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References

- [1] F. Berthelin, A. Vasseur, From kinetic equations to multidimensional isentropic gas dynamics before shocks, *SIAM J. Math. Anal.* 36(2005) 1807-1835.
- [2] M. Bostan, T. Goudon, High-electric-field limit for the Vlasov-Maxwell-Fokker-Planck system, preprint 2005.
- [3] Y. Brenier, Convergence of the Vlasov-Poisson system to the incompressible Euler equations, *Comm. Partial Differential Equations* 25(2000) 737-754.
- [4] Y. Brenier, N. Mauser, M. Puel, Incompressible Euler and e-MHD as scaling limits of the Vlasov-Maxwell system, *Commun. Math. Sci.* 1(2003) 437-447.
- [5] J.-Y. Chemin, *Fluides parfaits incompressibles*, Astérisque 230(1995).
- [6] P. Degond, M. Pulvirenti, Personal communication, 1994.
- [7] E. Frénod, E. Sonnendrücker, Homogenization of the Vlasov equation and of the Vlasov-Poisson system with strong external magnetic field, *Asymptotic Anal.* 18(1998) 193-213.

-
- [8] E. Frénod, E. Sonnendrücker, Long time behavior of the Vlasov equation with a strong external magnetic field, *Math. Models Methods Appl. Sci.* 10(2000) 539-553.
 - [9] R. Glassey, J. Schaeffer, On the 'one and one-half dimensional' relativistic Vlasov-Maxwell system, *Math. Methods Appl. Sci.* 13(1990) 169-179.
 - [10] R. Glassey, J. Schaeffer, The relativistic Vlasov-Maxwell system in two space dimensions, Part I and II, *Arch. Ration. Mech. Anal.* 141(1998) 331-354 and 355-374.
 - [11] F. Golse, L. Saint-Raymond, The Vlasov-Poisson system with strong magnetic field, *J. Math. Pures Appl.* 78(1999) 791-817.
 - [12] F. Golse, L. Saint-Raymond, The Vlasov-Poisson system with strong magnetic field in quasineutral regime, *Math. Models Methods Appl. Sci.* 13(2003) 661-714.
 - [13] T. Goudon, P.-E. Jabin, A. Vasseur, Hydrodynamic limits for the Vlasov-Navier-Stokes equations. Part II: Fine particles regime, *Indiana Univ. Math. J.* 53(2004) 1517-1536.
 - [14] T. Goudon, J. Nieto, F. Poupaud, J. Soler, Multidimensional high-field limit of the electrostatic Vlasov-Poisson-Fokker-Planck system, *J. Differential Equations* 213(2005) 418-442.
 - [15] H. Grad, The guiding center plasma, *Proc. Symp. Appl. Math. XVIII*, Amer. Math. Soc., Providence R.I.(1967) 162-248.
 - [16] E. Grenier, Pseudo-differential energy estimates of singular perturbations, *Comm. Pure Appl. Math.* 50(1997) 821-865.
 - [17] P.-E. Jabin, Large time concentrations for solutions to kinetic equations with energy dissipation, *Comm. Partial Differential Equations* 25(2000) 541-557.
 - [18] T. Kato, Nonstationary flows of viscous and ideal fluid in \mathbb{R}^3 , *J. Funct. Anal.* 9(1972) 296-305.
 - [19] E.M. Lifshitz, L.P. Pitaevskii, *Course of theoretical physics, Vol. 10*, translated from Russian by J.B. Sykes and R.N. Franklin, Pergamon Int. Library of Sci., Technology, Engineering and Social Studies, Pergamon Press, Oxford-Elmsford, N.Y. 1981.
 - [20] P.-L. Lions, *Mathematical topics in fluid mechanics, Vol.1 Incompressible models*, Oxford Lecture Series in Mathematics and its Applications, Oxford University Press, New-York, 1996.
 - [21] C. Marchioro, M. Pulvirenti, *Mathematical theory of incompressible nonviscous fluids*, Springer, New-York, 1994.
 - [22] L. Saint-Raymond, Convergence of solutions to the Boltzmann equation in the incompressible Euler limit, *Arch. Ration. Mech. Anal.* 166(2003) 47-80.

- [23] R. Temam, Local existence of C^∞ solutions of the Euler equations of incompressible perfect fluids, in *Turbulence and the Navier-Stokes equations*, Springer-Verlag, New-York, 1976 184-194.
- [24] H.T. Yau, Relative entropy and hydrodynamics of Ginzburg-Landau models, *Lett. Math. Phys.* 22(1991) 63-80.

Contents

1	Introduction	3
2	The two dimensional case	8
2.1	Analysis of the limit system	10
2.2	Evolution of the modulated energy	13
3	Other systems	21



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