

Ansgar Jüngel; Hailiang Li

Quantum Euler-Poisson systems: Existence of stationary states

*Archivum Mathematicum*, Vol. 40 (2004), No. 4, 435--456

Persistent URL: <http://dml.cz/dmlcz/107926>

## Terms of use:

© Masaryk University, 2004

Institute of Mathematics of the Academy of Sciences of the Czech Republic provides access to digitized documents strictly for personal use. Each copy of any part of this document must contain these *Terms of use*.



This paper has been digitized, optimized for electronic delivery and stamped with digital signature within the project *DML-CZ: The Czech Digital Mathematics Library* <http://project.dml.cz>

**QUANTUM EULER-POISSON SYSTEMS:  
EXISTENCE OF STATIONARY STATES**

ANSGAR JÜNGEL<sup>†</sup>, HAILIANG LI<sup>\*</sup>

ABSTRACT. A one-dimensional quantum Euler-Poisson system for semiconductors for the electron density and the electrostatic potential in bounded intervals is considered. The existence and uniqueness of strong solutions with positive electron density is shown for quite general (possibly non-convex or non-monotone) pressure-density functions under a “subsonic” condition, i.e. assuming sufficiently small current densities. The proof is based on a reformulation of the dispersive third-order equation for the electron density as a nonlinear elliptic fourth-order equation using an exponential transformation of variables.

1. INTRODUCTION

In 1927, Madelung gave a fluidynamical description of quantum systems governed by the Schrödinger equation for the wave function  $\psi$  :

$$\begin{aligned} i\varepsilon\partial_t\psi &= -\frac{\varepsilon^2}{2}\Delta\psi - \phi\psi \quad \text{in } \mathbb{R}^d \times (0, T), \\ \psi(\cdot, 0) &= \psi_0 \quad \text{in } \mathbb{R}^d, \end{aligned}$$

where  $T > 0$ ,  $d \geq 1$ ,  $\varepsilon > 0$  is the scaled Planck constant, and  $\phi = \phi(x, t)$  is some (given) potential. Separating the amplitude and phase of  $\psi = |\psi|\exp(iS/\varepsilon)$ , the particle density  $\rho = |\psi|^2$  and the particle current density  $j = \rho\nabla S$  for irrotational flows satisfy the so-called *Madelung equations* [26]

$$(1.1) \quad \partial_t\rho + \operatorname{div} j = 0,$$

$$(1.2) \quad \partial_t j + \operatorname{div} \left( \frac{j \otimes j}{\rho} \right) - \rho\nabla\phi - \frac{\varepsilon^2}{2}\rho\nabla \left( \frac{\Delta\sqrt{\rho}}{\sqrt{\rho}} \right) = 0 \quad \text{in } \mathbb{R}^d \times (0, T),$$

---

2000 *Mathematics Subject Classification*: 35J40, 35J60, 76Y05.

*Key words and phrases*: quantum hydrodynamics, existence and uniqueness of solutions, non-monotone pressure, semiconductors.

Received December 10, 2002.

where the  $i$ -th component of the convective term  $\operatorname{div}(j \otimes j/\rho)$  equals

$$\sum_{k=1}^d \frac{\partial}{\partial x_k} \left( \frac{j_i j_k}{\rho} \right).$$

The equations (1.1)–(1.2) can be interpreted as the pressureless Euler equations including the quantum Bohm potential

$$\frac{\varepsilon^2}{2} \frac{\Delta \sqrt{\rho}}{\sqrt{\rho}}.$$

They have been used for the modeling of superfluids like Helium II [23, 25].

Recently, Madelung-type equations have been derived for the modeling of quantum semiconductor devices, like resonant tunneling diodes, starting from the Wigner-Boltzmann equation [11] or from a mixed-state Schrödinger-Poisson system [14, 15]. There are several advantages of the fluidynamical description of quantum semiconductors. First, kinetic equations, like the Wigner equation, or Schrödinger systems are computationally very expensive, whereas for Euler-type equations efficient numerical algorithms are available [10, 30]. Second, the macroscopic description allows for a coupling of classical and quantum models. Indeed, setting the Planck constant  $\varepsilon$  in (1.2) equal to zero, we obtain the classical pressureless equations, so in both pictures, the same (macroscopic) variables can be used. Finally, as semiconductor devices are modeled in bounded domains, it is easier to find physically relevant boundary conditions for the macroscopic variables than for the Wigner function or for the wave function.

The Madelung-type equations derived by Gardner [11] and Gasser et al. [14] also include a pressure term and a momentum relaxation term taking into account interactions of the electrons in the semiconductor crystal, and are self-consistently coupled to the Poisson equation for the electrostatic potential  $\phi$ :

$$(1.3) \quad \partial_t \rho + \operatorname{div} j = 0,$$

$$(1.4) \quad \partial_t j + \operatorname{div} \left( \frac{j \otimes j}{\rho} \right) + \nabla p(\rho) - \rho \nabla \phi - \frac{\varepsilon^2}{2} \rho \nabla \left( \frac{\Delta \sqrt{\rho}}{\sqrt{\rho}} \right) = -\frac{j}{\tau},$$

$$(1.5) \quad \lambda^2 \Delta \phi = \rho - \mathcal{C}(x) \quad \text{in } \Omega \times (0, T),$$

where  $\Omega \subset \mathbb{R}^d$  is a bounded domain,  $\tau > 0$  is the (scaled) momentum relaxation time constant,  $\lambda > 0$  the (scaled) Debye length, and  $\mathcal{C}(x)$  is the doping concentration modeling the semiconductor device under consideration [18, 20, 28]. The pressure is assumed to depend only on the particle density and, like in classical fluid dynamics, often the expression

$$(1.6) \quad p(\rho) = \frac{T}{\gamma} \rho^\gamma, \quad \rho \geq 0,$$

with the temperature constant  $T > 0$  is employed [11, 19]. *Isothermal* fluids correspond to  $\gamma = 1$ , *isentropic* fluids to  $\gamma > 1$ . Notice that the particle temperature is  $T(\rho) = T\rho^{\gamma-1}$ .

The equations (1.3)–(1.5) are referred to as the *quantum Euler-Poisson system* or as the quantum hydrodynamic model.

In this paper we study the stationary system (1.3)–(1.5) in one space dimension with  $\lambda = 1$ :

$$(1.7) \quad j_0 = \text{const.},$$

$$(1.8) \quad \left( \frac{j_0^2}{\rho} + p(\rho) \right)_x - \rho \phi_x - \frac{\varepsilon^2}{2} \rho \left( \frac{(\sqrt{\rho})_{xx}}{\sqrt{\rho}} \right)_x = -\frac{j_0}{\tau},$$

$$(1.9) \quad \phi_{xx} = \rho - \mathcal{C}(x) \quad \text{in } \Omega = (0, 1)$$

subject to the boundary conditions

$$(1.10) \quad \rho(0) = \rho_1, \quad \rho(1) = \rho_2, \quad \rho_x(0) = \rho_x(1) = 0,$$

$$(1.11) \quad \phi(0) = 0, \quad \phi(1) = \Phi_0,$$

where  $\rho_1, \rho_2 > 0$  and  $\Phi_0 \in \mathbb{R}$ . In this formulation, the electron current density is a given constant. From the equations, the applied voltage  $U$  can be computed by  $U = \phi(1) - \phi(0)$ .

As the momentum equation (1.8) is of third order, the mathematical analysis of the above system of equations is quite difficult. In fact, without the third-order quantum term, the above equations represent the Euler-Poisson system of gas dynamics for which only partial existence results (in several space dimensions) are available (see, e.g., [5, 9] for several space dimensions and [6] for one space dimension).

Therefore, we can only expect partial results for the hydrodynamic equations *including* the third-order quantum term which makes the problem even more difficult. In the following, we describe some mathematical techniques which have been successfully applied to the system (1.7)–(1.9) to prove the existence (and uniqueness) of solutions.

In the literature, there exist essentially two ideas in dealing with the nonlinear third-order equation (see also [14]). One idea consists in reducing the momentum equation (1.8) to a *second-order* equation. The second idea is to differentiate (1.8) once and to obtain a *fourth-order* equation.

The first idea has been used in [8, 7, 19, 31]. The existence of solutions to (1.7)–(1.9) has been shown for sufficiently small  $j_0 > 0$ , using nonlinear boundary conditions for  $\sqrt{\rho}_{xx}$  or Dirichlet data for the velocity potential. The pressure function is assumed to be a monotone function of the density.

The second idea has been employed in [16] in order to prove the existence of solutions to (1.7)–(1.9), again for sufficiently small  $j_0 > 0$ . In that work, the boundary conditions (1.10)–(1.11) have been used, but the pressure has been assumed to be linear:  $p(\rho) = \rho$ . The main idea in [16] was to write the density in exponential form:  $n = e^u$  and to derive uniform  $H^1$  bounds for  $u$  which, by Sobolev embedding, yields  $L^\infty$  bounds for  $u$  and hence a positive lower bound for  $n = e^u$ .

The main aim of this paper is to generalize the results of [16] to *general* pressure functions. Compared to the results in [8], we use different boundary conditions

and more general pressure functions. Moreover, the technique of proof is different. Compared to [16], we allow for more general pressure functions, in particular also *non-convex* or *non-monotone* pressure-density relations.

We mention some related results on the stationary quantum Euler-Poisson system. The semi-classical limit  $\varepsilon \rightarrow 0$  in the case of thermal equilibrium  $j_0 = 0$  and in the case  $j_0 > 0$  has been studied in [12, 29] and [16], respectively (also see [13]). For results on the limit problem  $\varepsilon = 0$  (Euler-Poisson system) we refer to the review paper [24]. The local existence of strong solutions to the *transient* quantum Euler-Poisson model has been shown in [21]. The global existence of “small” solutions to the transient model and its asymptotic behavior for large times will be studied in our forthcoming work [22] based on the results of this paper for the steady state.

In all cited papers, the existence of (strong) steady-state solutions to the quantum hydrodynamic equations is shown for sufficiently small current densities  $j_0 > 0$ . In fact, in the case of the nonlinear boundary conditions assumed in [8], the *non-existence* of weak solutions to the quantum Euler-Poisson system for sufficiently large  $j_0 > 0$  has been proved. We also need the smallness condition on  $|j_0|$  to prove the existence of solutions to (1.7)–(1.11).

In order to explain our main results in detail, we rewrite the equation for the electron density (1.8) as a nonlinear elliptic *fourth-order* equation and write the density in exponential form. Writing

$$\frac{\varepsilon^2}{2} \rho \left( \frac{(\sqrt{\rho})_{xx}}{\sqrt{\rho}} \right)_x = \frac{\varepsilon^2}{4} (\rho(\ln \rho)_{xx})_x,$$

dividing (1.8) by  $\rho > 0$ , differentiating with respect to  $x$  and using (1.7) and (1.9) to remove the electrostatic potential from the equation, we obtain

$$\left( \frac{p'(\rho)\rho_x}{\rho} - \frac{j_0^2}{\rho^3} \rho_x \right)_x - (\rho - \mathcal{C}(x)) - \frac{\varepsilon^2}{4} (\rho^{-1}(\rho(\ln \rho)_{xx})_x)_x = - \left( \frac{j_0}{\tau \rho} \right)_x.$$

It is convenient to introduce the new variable  $u = \ln \rho$ . Then the above equation can be written as

$$(1.12) \quad \frac{\varepsilon^2}{4} \left( u_{xx} + \frac{1}{2} u_x^2 \right)_{xx} - ((p'(e^u) - j_0^2 e^{-2u})u_x)_x + e^u - \mathcal{C}(x) = \frac{j_0}{\tau} (e^{-u})_x.$$

The boundary conditions (1.10) transform to

$$(1.13) \quad u(0) = u_1, \quad u(1) = u_2, \quad u_x(0) = u_x(1) = 0,$$

where  $u_1 = \ln \rho_1$ ,  $u_2 = \ln \rho_2$ .

The electrostatic potential can be computed from the formulae

$$(1.14) \quad \phi(x) = \Phi_0 + \int_0^1 G(x, y)(e^{u(y)} - \mathcal{C}(y)) dy,$$

where the Green’s function  $G(x, y)$  is defined by

$$G(x, y) = \begin{cases} x(1 - y), & x < y, \\ y(1 - x), & x > y. \end{cases}$$

The advantage of the above formulation is that bounded solutions  $u \in L^\infty(0, 1)$  define *positive* densities  $\rho = e^u$  and in this case, both formulations (1.8)–(1.9) and (1.12)–(1.14) are equivalent. Notice that for third-order or fourth-order equations, no maximum principle is available such that other methods for proving the positivity of the variables have to be devised. Here we use the exponential transformation of variables combined with Sobolev embeddings as in [8, 16].

Assume that

$$(1.15) \quad \varepsilon, \tau, \rho_1, \rho_2 > 0, \quad \Phi_0, j_0 \in \mathbb{R}, \quad C \in L^2(0, 1).$$

Then our main results are as follows:

1. Suppose that the pressure function is given by (1.6) for  $\gamma > 0$ . Then there exist constants  $J_0, \gamma_1 > 0$  such that if  $|j_0| \leq J_0$  and  $|\gamma - 1| \leq \gamma_1$ , there exists a unique strong solution  $u, \phi \in H^4(0, 1)$  to (1.12)–(1.14). Since  $u \in L^\infty(0, 1)$ , we have  $\rho = e^u > 0$  in  $(0, 1)$ , and  $\rho, \phi \in H^4(0, 1)$  is a solution of (1.7)–(1.11). The constant  $j_1$  can be given explicitly (see section 2).
2. Suppose that  $p \in C^3(0, \infty)$ , that there exists a function  $\mathcal{A} \in H^2(0, 1)$  such that

$$\mathcal{A} > 0 \quad \text{in} \quad (0, 1), \quad \mathcal{A}(0) = \rho_1, \quad \mathcal{A}(1) = \rho_2, \quad \mathcal{A}_x(0) = \mathcal{A}_x(1) = 0$$

and that there is a set  $E \subset [0, 1]$  such that

$$(1.16) \quad p'(\mathcal{A}) - \frac{j_0^2}{\mathcal{A}^2} \begin{cases} \leq 0, & x \in E, \\ > 0, & x \in [0, 1] \setminus E. \end{cases}$$

Then if  $|j_0|$  is small enough, there exists a unique strong solution  $u, \phi \in H^4(0, 1)$  to (1.12)–(1.14).

Notice that we allow for *non-convex* pressure functions (1.6) with  $\gamma < 1$  and for *non-monotone* pressures satisfying (1.16). This means that the left part of (1.2) may be not hyperbolic. The assumption (1.16) implies that the interval under consideration may consist of subsonic, transonic and supersonic regions in the classical sense [3]. To guarantee the well-posedness of strong solutions, we assume a “subsonic” condition.

Finally, we notice that our estimates allow to perform the semi-classical limit  $\varepsilon \rightarrow 0$  in (1.12)–(1.14) by employing the same techniques as in [16] (also see [7] and Remark 3.4).

The paper is organized as follows. In section 2 our first main result is formulated and proved. The second main result is shown in section 3.

**Notation.** The Lebesgue space of square integrable functions with the norm  $\|\cdot\|$  is denoted by  $L^2(0, 1)$ , and  $H^k(0, 1)$  or simply  $H^k$  denotes the usual Sobolev space of functions  $f$  satisfying  $\partial_x^i f \in L^2(0, 1)$ ,  $0 \leq i \leq k$ , with the norm  $\|\cdot\|_k$ . In particular,  $\|\cdot\|_0 = \|\cdot\|$ .

2. PRESSURE FUNCTIONS SATISFYING (1.6)

In this section, we consider the steady-state solutions to the BVP (1.7)–(1.11) when the pressure-density relation satisfies the  $\gamma$ -law (1.6).

**Theorem 2.1.** *Assume that (1.6) and (1.15) hold. Let  $\kappa \in (0, 1)$ . Then there exist two constants  $\gamma_0 > 0$  and  $K(\kappa) > 0$  such that if*

$$(2.1) \quad |j_0| \leq e^{-K(\kappa)} \kappa \sqrt{Te^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2} \quad \text{and} \quad |\gamma - 1| \leq \gamma_0,$$

then there exists a solution  $u \in H^4$  to the BVP (1.12)–(1.13) satisfying

$$(2.2) \quad \frac{1}{2}\varepsilon \|u_{xx}\| + \sqrt{Te^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2} \|u_x\| \leq K_0,$$

$$(2.3) \quad |u(x)| \leq K(\kappa),$$

where  $K_0$  is defined by (2.21) and  $(\gamma, K(\kappa))$  is the unique solution to (2.22).

Furthermore, there are  $J_0, \varepsilon_0, \gamma_1 > 0$  such that if  $|j_0| \leq J_0$  and  $|\gamma - 1| \leq \gamma_1$ , the solution  $u$  is unique for any  $\varepsilon \in (0, \varepsilon_0]$ .

**Proof.** *Step 1. A-priori estimates.* Assume that  $u \in H^2$  is a weak solution of the boundary-value problem (BVP) (1.12)–(1.13) satisfying a-priorily that

$$(2.4) \quad -K(\kappa) \leq u \leq K(\kappa).$$

Following [1] we introduce a function  $u_D \in C^2([0, 1])$  satisfying

$$(2.5) \quad u_D(0) = u_1, \quad u_D(1) = u_2, \quad u_{D,x}(0) = u_{D,x}(1) = 0,$$

with piecewise linear second order derivative

$$(2.6) \quad u_{D,xx}(x) = \begin{cases} \frac{4\zeta}{\mu^2(1-\mu)}x, & x \in [0, \frac{\mu}{2}), \\ \frac{4\zeta}{\mu^2(1-\mu)}(\mu - x), & x \in [\frac{\mu}{2}, \mu], \\ 0, & x \in (\mu, \frac{1}{2}], \end{cases}$$

and  $u_{D,xx}(x) = -u_{D,xx}(1 - x)$  for  $x \in (\frac{1}{2}, 1]$ , where  $\zeta = |u_1 - u_2|$  and  $\mu \in (0, \frac{1}{2})$ . Elementary computations show that

$$(2.7) \quad \int_0^{1/2} x|u_{D,xx}(x)| dx + \int_{1/2}^1 (1-x)|u_{D,xx}(x)| dx = \frac{\mu\zeta}{1-\mu},$$

$$(2.8) \quad \int_0^1 |u_{D,xx}(x)|^2 dx = \frac{8\zeta^2}{3\mu(1-\mu)},$$

$$(2.9) \quad \int_0^1 |u_{D,x}(x)| dx = \zeta,$$

$$(2.10) \quad \int_0^1 |u_{D,x}(x)|^2 dx = \frac{\zeta^2(23\mu + 1)}{30(1-\mu)^2} \leq \frac{1}{2}\zeta^2.$$

Use  $u - u_D \in H_0^2$  as an admissible test function in the weak formulation of (1.12) to obtain

$$\begin{aligned}
 & \frac{1}{4}\varepsilon^2 \int_0^1 (u_{xx} + \frac{1}{2}u_x^2)(u - u_D)_{xx} dx \\
 & \quad + \int_0^1 (Te^{(\gamma-1)u} - j_0^2 e^{-2u})u_x(u - u_D)_x dx \\
 & \quad + \int_0^1 (e^u - e^{u_D})(u - u_D) dx - \int_0^1 (e^{u_D} - \mathcal{C})(u - u_D) dx \\
 (2.11) \quad & = -\frac{j}{\tau} \int_0^1 e^{-u}(u - u_D)_x dx,
 \end{aligned}$$

By Cauchy's inequality, (2.4) and (2.11), it follows for  $\eta \in (0, 1)$

$$\begin{aligned}
 & \frac{1}{4}\varepsilon^2 \left(1 - \frac{\eta}{2}\right) \int_0^1 u_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 u_x^2 u_{xx} dx \\
 & \quad + \int_0^1 \left[\left(1 - \frac{\eta}{2}\right) Te^{-|\gamma-1|K(\kappa)} - (1 + \eta)j_0^2 e^{2K(\kappa)}\right] u_x^2 dx \\
 & \quad + \int_0^1 (e^u - e^{u_D})(u - u_D) dx \\
 & \leq \frac{1}{8} \frac{\varepsilon^2}{\eta} \int_0^1 |u_{D,xx}|^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 u_x^2 u_{D,xx} dx + \frac{j_0^2}{4\eta} \int_0^1 e^{2K(\kappa)} |u_{D,x}|^2 dx \\
 & \quad + \frac{1}{2\eta} Te^{2|\gamma-1|K(\kappa)} \int_0^1 |u_{D,x}|^2 dx + \int_0^1 (e^{u_D} - \mathcal{C})(u - u_D) dx \\
 (2.12) \quad & - \frac{j_0}{\tau} \int_0^1 e^{-u}(u - u_D)_x dx,
 \end{aligned}$$

Using the boundary condition (1.13), and applying (2.5) and (2.7)–(2.8), we have as the proof of Lemma 2.1 in [16]:

$$\begin{aligned}
 & \int_0^1 u_x^2 u_{xx} dx = 0, \quad \frac{1}{8} \frac{\varepsilon^2}{\eta} \int_0^1 |u_{D,xx}|^2 dx \leq \frac{\zeta^2 \varepsilon^2}{3\mu\eta}, \\
 (2.13) \quad & \frac{1}{8}\varepsilon^2 \int_0^1 u_x^2 u_{D,xx} dx \leq \frac{\varepsilon^2 \eta}{16} \|u_{xx}\|^2,
 \end{aligned}$$

where we have used the Poincaré inequality and chosen  $\mu = \min(1/2, \eta/2\zeta)$ . From (2.5) and (2.10) follows

$$(2.14) \quad \frac{j_0^2}{4\eta} \int_0^1 e^{2K(\kappa)} |u_{D,x}|^2 dx \leq \frac{\zeta^2}{8\eta} j_0^2 e^{2K(\kappa)} \leq \frac{\zeta^2}{8\eta} \left( Te^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2 \right)$$

and

$$(2.15) \quad \frac{1}{2\eta} Te^{2|\gamma-1|K(\kappa)} \int_0^1 |u_{D,x}|^2 dx \leq \frac{\zeta^2}{4\eta} Te^{2|\gamma-1|K(\kappa)}.$$



Since

$$\|u - u_D\| \leq \frac{1}{2} \|(u - u_D)_x\|,$$

we have by Cauchy's inequality

$$\begin{aligned} & \int_0^1 (e^{u_D} - \mathcal{C})(u - u_D) \, dx \\ & \leq \frac{\eta}{8} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right) \|(u - u_D)_x\|^2 \\ & \quad + \frac{1}{2\eta} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right)^{-1} \|e^{u_D} - \mathcal{C}\|^2 \\ & \leq \frac{\eta}{8} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right) \|u_x\|^2 + \frac{\eta \zeta^2}{16} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right) \\ (2.16) \quad & + \frac{1}{2\eta} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right)^{-1} \|e^{u_D} - \mathcal{C}\|^2. \end{aligned}$$

Using (2.1) and (2.9), the last term in (2.12) can be estimated as

$$\begin{aligned} & - \frac{j_0}{\tau} \int_0^1 e^{-u} (u - u_D)_x \, dx \\ & \leq \frac{|j_0|}{\tau} e^{K(\kappa)} \int_0^1 (|u_x| + |u_{D,x}|) \, dx \\ (2.17) \quad & \leq \frac{\eta}{4} \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right) \|u_x\|^2 + \frac{1}{\eta \tau^2} + \frac{\zeta \kappa}{\tau} \sqrt{T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2}. \end{aligned}$$

Setting  $\eta = (1 - \kappa)/[2(1 + \kappa)]$  in (2.13)–(2.17) and substituting  $\eta$  into (2.12), we have, estimating as in [16],

$$(2.18) \quad \frac{1}{4} \varepsilon^2 \|u_{xx}\|^2 + \left( T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 \right) \|u_x\|^2 \leq a_c^2 \frac{(1 + \kappa)^2}{(1 - \kappa)^2} K_1,$$

where  $K_1$  is given by

$$(2.19) \quad K_1 = 1 + T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2 + T e^{2|\gamma-1|K(\kappa)} + T^{-1} e^{|\gamma-1|K(\kappa)},$$

and  $a_c$  a generic positive constant which only depends on  $\zeta$ ,  $T$ , and  $\tau$ . From (2.18) follows

$$(2.20) \quad \frac{1}{2} \varepsilon \|u_{xx}\| + \sqrt{T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2} \|u_x\| \leq K_0,$$

where

$$(2.21) \quad K_0 = a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{K_1}.$$

Now, consider the equation for  $(\gamma, K(\kappa))$

$$(2.22) \quad K(\kappa) - a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{K_2(K(\kappa), \gamma)} - |u_1| = 0,$$

where

$$K_2(K(\kappa), \gamma) = 1 + T^{-1}e^{|\gamma-1|K(\kappa)} + e^{3|\gamma-1|K(\kappa)} + T^{-2}e^{2|\gamma-1|K(\kappa)}.$$

It has a solution

$$(\gamma, K(\kappa)) = (1, |u_1| + a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{2 + T^{-1} + T^{-2}}).$$

By the implicit function theorem, there exists a  $\gamma_0 > 0$  such that for  $|\gamma - 1| < \gamma_0$ , the equation (2.22) has a solution  $(\gamma, K(\kappa))$ .

Therefore, in view of (2.19)–(2.22), we obtain

$$\begin{aligned} |u(x)| &\leq |u_1| + \|u_x\| \\ &= |u_1| + \frac{K_0}{\sqrt{T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2}} \\ &= |u_1| + a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{\frac{1 + T e^{-|\gamma-1|K} + \frac{1}{2}\varepsilon^2 + T e^{2|\gamma-1|K(\kappa)} + T^{-1} e^{|\gamma-1|K(\kappa)}}{T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2}} \\ &\leq |u_1| + a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{K_2(K(\kappa), \gamma)} \\ (2.23) \quad &= K(\kappa). \end{aligned}$$

*Step 2. Existence.* We apply the Leray-Schauder fixed point theorem to prove the existence of strong solutions. Let  $v \in X := C^{0,1}([0, 1])$ . Consider the linear BVP

$$\begin{aligned} \frac{1}{4}\varepsilon^2 \left( u_{xx} + \frac{1}{2}\lambda v_x^2 \right)_{xx} - \left( T e^{(\gamma-1)v_K} u_x - \lambda j_0^2 e^{2v_K} u_x \right)_x \\ + \lambda \left( \frac{e^v - 1}{v} u + 1 - \mathcal{C} \right) = \lambda \frac{j_0}{\tau} (e^{-v_K})_x, \\ u(0) = \lambda u_1, \quad u(1) = \lambda u_2, \quad u_x(0) = u_x(1) = 0, \end{aligned}$$

where  $v_K = \min\{K(\kappa), \max\{-K(\kappa), v\}\}$  and  $\lambda \in [0, 1]$ . Define the bilinear form

$$a(u, \psi) = \int_0^1 \left( \frac{1}{4}\varepsilon^2 u_{xx} \psi_{xx} + T e^{(\gamma-1)v_K} u_x \psi_x + \lambda \frac{e^v - 1}{v} u \psi \right) dx,$$

for  $u, \psi \in H^2$  and the functional

$$F(\psi) = \int_0^1 \left( -\frac{1}{8}\varepsilon^2 \lambda v_x^2 \psi_{xx} + \lambda j_0^2 e^{-2v_K} v_x \psi_x + \lambda (\mathcal{C} - 1) \psi - \lambda \frac{j_0}{\tau} e^{-v_K} \psi \right) dx.$$

Since  $X \hookrightarrow W^{1,\infty}$ ,  $a(\cdot, \cdot)$  is continuous and coercive in  $H^2$ , and  $F$  is linear and continuous in  $H^2$ , the Lax-Milgram theorem yields the existence of a solution  $u \in H^2$ . This means that the map  $\mathcal{S} : X \times [0, 1] \rightarrow X$ ,  $(v, \lambda) \mapsto u$  is well defined. Moreover, it is not difficult to see that  $\mathcal{S}$  is continuous and compact.

Since  $\mathcal{S}(v, 0) = 0$  for all  $v \in H^2$ , and by similar estimates as in Step 1 with  $e^u$  replaced by  $e^{v\kappa}$ , we can verify that it holds for all  $\lambda \in [0, 1]$ ,

$$(2.24) \quad \|u\|_X \leq c_1,$$

where  $c_1 > 0$  is a constant independent of  $u$  and  $\lambda$ . Then the existence of  $u \in H^2$  follows from the Leray-Schauder fixed point theorem. It is not difficult to prove that indeed  $u \in H^4$  (see [16] for details).

*Step 3. Uniqueness.* Let  $u, v \in H_0^2$  be two weak solutions of the BVP (1.12)–(1.13), which satisfy (2.20). Using  $u - v \in H_0^2$  as an admissible test function in the weak formulation derived for  $u - v$ , we obtain

$$\begin{aligned} & \frac{1}{4}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x + v_x)(u - v)_x(u - v)_{xx} dx \\ & \quad + T e^{(\gamma-1)u} \int_0^1 (u - v)_x^2 dx \\ & = - \int_0^1 \left( T e^{(\gamma-1)u} - T e^{(\gamma-1)v} \right) v_x (u - v)_x dx \\ & \quad + \frac{1}{2} \int_0^1 (j_0^2 e^{-2u} - j_0^2 e^{-2v})(u - v)_{xx} dx \\ & \quad - \int_0^1 (e^u - e^v)(u - v) dx - \frac{j_0}{\tau} \int_0^1 (e^{-u} - e^{-v})(u - v)_x dx \\ & \leq \int_0^1 T e^{|\gamma-1|K(\kappa)} \left( e^{(\gamma-1)u} - e^{(\gamma-1)v} \right)^2 v_x^2 dx \\ & \quad + \frac{1}{8}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{2\varepsilon^2} j_0^4 \int_0^1 (e^{-2u} - e^{-2v})^2 dx \\ & \quad - e^{-K(\kappa)} \int_0^1 (u - v)^2 dx + \frac{1}{2} T e^{-|\gamma-1|K(\kappa)} \int_0^1 (u - v)_x^2 dx \\ & \quad + \frac{j_0^2}{\tau^2 T} e^{|\gamma-1|K(\kappa)} \int_0^1 (e^{-u} - e^{-v})^2 dx, \end{aligned}$$

which implies

$$\begin{aligned} & \frac{1}{8}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x + v_x)(u - v)_x(u - v)_{xx} dx \\ & \quad + \frac{1}{2} T e^{-|\gamma-1|K} \int_0^1 (u - v)_x^2 dx \\ & \leq \left( \frac{1}{2\varepsilon^2} j_0^4 e^{4K(\kappa)} + \frac{j_0^2}{\tau^2 T} e^{(|\gamma-1|+2)K} + T(\gamma - 1)^2 e^{3|\gamma-1|K(\kappa)} |v_x|_{L^\infty}^2 \right) \\ (2.25) \quad & \times \int_0^1 (u - v)^2 dx - e^{-K(\kappa)} \int_0^1 (u - v)^2 dx. \end{aligned}$$

From (1.13) and (2.20), we have by Hölder’s inequality

$$\begin{aligned} |u_x(x)| &\leq \sqrt{2\|u_{xx}\| \cdot \|u_x\|} \leq \frac{1}{2}\theta\|u_{xx}\| + \theta^{-1}\|u_x\| \\ &\leq \max\left\{\frac{\theta}{\varepsilon}, \frac{1}{\theta T^{\frac{1}{2}}e^{-\frac{1}{2}|\gamma-1|K(\kappa)}}\right\} \tilde{K}_0, \end{aligned}$$

where

$$\tilde{K}_0 = a_c \frac{1 + \kappa}{1 - \kappa} \sqrt{K_1}.$$

Choose

$$\theta = \frac{T^{\frac{1}{2}}e^{-\frac{1}{2}|\gamma-1|K(\kappa)}}{\tilde{K}_0} \quad \text{and} \quad 0 < \varepsilon \leq \varepsilon_0 =: \frac{T^{\frac{3}{2}}e^{-\frac{3}{2}|\gamma-1|K(\kappa)}}{\tilde{K}_0^2}$$

to obtain

$$\frac{\theta}{\varepsilon} \geq \frac{1}{\theta T^{\frac{1}{2}}e^{-\frac{1}{2}|\gamma-1|K(\kappa)}}$$

and

$$(2.26) \quad |u_x(x)| \leq \frac{\theta}{\varepsilon} \tilde{K}_0 \leq \frac{1}{\varepsilon} T^{\frac{1}{2}}e^{-\frac{1}{2}|\gamma-1|K(\kappa)}.$$

From this estimate and Cauchy’s inequality follows that the left-hand side of (2.25) is bounded from below by

$$\begin{aligned} &\frac{1}{8}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x + v_x)(u - v)_x(u - v)_{xx} dx \\ &+ \frac{1}{2} \int_0^1 T e^{-|\gamma-1|K(\kappa)} (u - v)_x^2 dx \\ &\geq \frac{1}{16}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{4} T e^{-|\gamma-1|K(\kappa)} \int_0^1 (u - v)_x^2 dx. \end{aligned}$$

On the other hand, by the implicit function theorem, there is a  $\gamma_1 \leq \gamma_0$  such that for  $|\gamma - 1| < \gamma_1$  it holds

$$(2.27) \quad |\gamma - 1|^2 e^{(2|\gamma-1|+1)K(\kappa)} \leq \frac{1}{8} T^{-2} \varepsilon^2,$$

which implies

$$(2.28) \quad \frac{2}{\varepsilon^2} T^2 (\gamma - 1)^2 e^{2|\gamma-1|K(\kappa)} \leq \frac{1}{4} e^{-K(\kappa)}.$$

Thus, there exists  $J_0$  such that if

$$(2.29) \quad j_0^2 \leq J_0^2 =: \min\left\{e^{-2K(\kappa)} \kappa \left(T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2}\varepsilon^2\right), \frac{\sqrt{2}}{2}\varepsilon e^{-\frac{5}{2}K(\kappa)}, \frac{1}{4} T \tau^2 e^{-(|\gamma-1|+3)K(\kappa)}\right\},$$

then the first integral on the right-hand side of (2.25) is bounded, in view of (2.26) and (2.28), by

$$\begin{aligned}
 & \left( \frac{1}{2\varepsilon^2} j_0^4 e^{4K(\kappa)} + \frac{j_0^2}{\tau^2 T} e^{(|\gamma-1|+2)K(\kappa)} + T(\gamma-1)^2 e^{3|\gamma-1|K(\kappa)} |v_x|_{L^\infty}^2 - e^{-K(\kappa)} \right) \\
 & \quad \times \int_0^1 (u-v)^2 dx \\
 & \leq \left( \frac{1}{2\varepsilon^2} j_0^4 e^{4K(\kappa)} + \frac{j_0^2}{\tau^2 T} e^{(|\gamma-1|+2)K(\kappa)} + \frac{2}{\varepsilon^2} T^2 (\gamma-1)^2 e^{2|\gamma-1|K(\kappa)} - e^{-K(\kappa)} \right) \\
 & \quad \times \int_0^1 (u-v)^2 dx \\
 (2.30) \quad & \leq -\frac{1}{4} e^{-K(\kappa)} \int_0^1 (u-v)^2 dx.
 \end{aligned}$$

Therefore, the weak solution is unique if both (2.27) and (2.29) holds. The proof of Theorem 2.1 is complete. □

As in [16], we can conclude from Theorem 2.1 the following result.

**Theorem 2.2.** *Assume that (1.6) and (1.15) hold. Then there exist two constants  $\gamma_0 > 0$  and  $K(\kappa) > 0$  ( $\kappa \in (0, 1)$ ) such that if*

$$(2.31) \quad |j_0| \leq e^{-K(\kappa)} \kappa \sqrt{T e^{-|\gamma-1|K(\kappa)} + \frac{1}{2} \varepsilon^2} \quad \text{and} \quad |\gamma-1| \leq \gamma_0,$$

then there is a solution  $(\rho, \phi) \in H^4 \times H^2$  to the BVP (1.7)–(1.11) such that

$$(2.32) \quad \rho \geq \bar{\rho} =: e^{-K(\kappa)} > 0,$$

where  $(\gamma, K(\kappa))$  solves the equation (2.22). Moreover, if  $|j_0|$ ,  $\varepsilon$ , and  $|\gamma-1|$  are small enough, the solution is unique.

### 3. PRESSURE FUNCTIONS SATISFYING (1.16)

In this section, we consider the BVP (1.12)–(1.13) (and (1.7)–(1.11)) with pressure functions satisfying the condition (1.16).

Set  $u_D := \ln \mathcal{A}$ . Then  $e^{u_D} = \mathcal{A}$ . We have the following theorem.

**Theorem 3.1.** *Assume that (1.15) and (1.16) holds and that  $p \in C^3(0, \infty)$ . For  $\kappa \in (0, 1)$  assume that it holds*

$$(3.1) \quad \min_{x \in [0,1]} \mathcal{A}(x)^2 \left( \frac{1}{4} \kappa \varepsilon^2 + p'(\mathcal{A}(x)) \right) > j_0^2.$$

Then the BVP (1.12)–(1.13) has a unique solution  $u \in H^4$  provided that  $\|\mathcal{A}'\|_1 + \|\mathcal{A} - \mathcal{C}\|$  is sufficiently small. Moreover, it holds

$$(3.2) \quad \mathcal{A}_* \|u - u_D\|^2 + A_0 \|u_{xx}\|^2 + \varepsilon^4 \|(u_{xxx}, u_{xxxx})\|^2 + \int_{I \setminus E} \left( p'(\mathcal{A}^2) - \frac{j_0^2}{\mathcal{A}^4} \right) (u - u_D)_x^2 dx \leq K_c \delta_0,$$

where

$$(3.3) \quad \mathcal{A}_* = \min_{x \in [0,1]} \mathcal{A}(x), \quad \delta_0 = \|\mathcal{A}'\|_1 + \|\mathcal{A} - \mathcal{C}\|,$$

$$(3.4) \quad A_0 = \min_{x \in [0,1]} \left( \frac{1}{4} \kappa \varepsilon^2 + p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2} \right) > 0,$$

and  $K_c > 0$  is a constant depending on  $\mathcal{A}$ ,  $\tau$  and  $j_0$ .

**Remark 3.2.** (1) We call the main assumption (3.1) a “subsonic” condition for the quantum Euler-Poisson system. When  $\varepsilon = 0$ , the assumption (3.1) is exactly the subsonic condition for the classical hydrodynamic model [4].

(2) One can verify that the assumption (3.1) can be replaced by

$$(3.5) \quad \frac{1}{4} \kappa \varepsilon^2 + \text{meas}(E) \min_{x \in E} (p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2}) > 0, \quad \kappa \in (0, 1),$$

in order to obtain the existence and uniqueness of strong solutions. Here, we recall that the region  $E \subset [0, 1]$  is defined such that it holds  $p'(\mathcal{A})\mathcal{A}^2 - j_0^2 < 0$  in  $E$ .

**Proof.** We prove Theorem 3.1 by the same steps as Theorem 2.1.

*Step 1. The a-priori estimates.* Let  $u \in H^4$  be a solution of the BVP (1.12)–(1.13) satisfying

$$(3.6) \quad u_D - \delta_1 \leq u \leq u_D + \delta_1,$$

where  $\delta_1 > 0$  is chosen such that

$$(3.7) \quad \frac{4}{5} \mathcal{A}_* \leq e^{-\delta_1} e^{u_D} \leq e^u \leq e^{\delta_1} e^{u_D} \leq \frac{5}{4} \mathcal{A}^*,$$

$$(3.8) \quad \max_{x \in [0,1]} (|p''(\mathcal{A})| \mathcal{A} + 2j_0^2 \mathcal{A}^{-2}) \delta_1 \leq \frac{2}{9} (1 + \theta) A_0,$$

$$(3.9) \quad \max_{\ln \frac{16}{25} \mathcal{A}_* \leq y \leq \ln \frac{25}{16} \mathcal{A}^*} (|p'''(e^y)| e^{2y} + |p''(e^y)| e^y + 4j_0^2 e^{-2y}) \delta_1^2 \leq \frac{4}{9} (1 + \theta) A_0.$$

where  $\mathcal{A}^* = \max_{x \in [0,1]} \mathcal{A}(x)$ ,  $\mathcal{A}_* = \min_{x \in [0,1]} \mathcal{A}(x)$ , and  $\theta = \frac{1-\kappa}{1+\kappa}$  (then  $\kappa = \frac{1-\theta}{1+\theta}$ ).

Assume that  $\delta_0 = \|\mathcal{A} - \mathcal{C}\| + \|(\mathcal{A}_x, \mathcal{A}_{xx})\|$  is so small that it holds

$$(3.10) \quad |u_{D,x}|_\infty + |u_{D,xx}|_\infty \leq \theta,$$

where  $|\cdot|_\infty$  denotes the  $L^\infty$  norm.

Taking  $(u - u_D) \in H_0^2$  as an admissible test function in the weak formulation of (1.12), we have, by Cauchy’s inequality and (3.7), that

$$\frac{1}{4} \left( 1 - \frac{1}{2} \theta - \frac{1}{2} |u_{D,xx}|_\infty \right) \varepsilon^2 \int_0^1 u_{xx}^2 dx$$

$$\begin{aligned}
 & + \int_0^1 (p'(e^u) - j_0^2 e^{-2u}) u_x (u - u_D)_x dx + \frac{1}{2} \mathcal{A}_* \int_0^1 (u - u_D)^2 dx \\
 (3.11) \quad & \leq \frac{1}{\mathcal{A}_*} \int_0^1 (\mathcal{A} - \mathcal{C})^2 dx + \frac{1}{8\theta} \varepsilon^2 \int_0^1 |u_{D,xx}|^2 dx - \frac{j_0}{\tau} \int_0^1 e^{-u} (u - u_D)_x dx,
 \end{aligned}$$

where we have used the facts that

$$(3.12) \quad \int_0^1 u_x^2 u_{xx} dx = 0, \quad \int_0^1 u_x^2 dx \leq \int_0^1 u_{xx}^2 dx.$$

The last term in (3.11) can be estimated as

$$\begin{aligned}
 -\frac{j_0}{\tau} \int_0^1 e^{-u} (u - u_D)_x dx & = \frac{j_0}{\tau} e^{-u_D} e^{-(u-u_D)} \Big|_{x=0}^1 + \frac{j_0}{\tau} \int_0^1 e^{-u} u_{D,x} dx \\
 (3.13) \quad & = -\frac{j_0}{\tau} \int_0^1 (e^{-u_D} - e^{-u}) u_{D,x} dx \leq K_c \delta_0,
 \end{aligned}$$

where here and in the following  $K_c > 0$  is a generic constant depending on  $\mathcal{A}$ ,  $\tau$  and  $j_0$ .

By Taylor's expansion and Cauchy's inequality, the second term on the left-hand side of (3.11) can be estimated as

$$\begin{aligned}
 & \int_0^1 (p'(e^u) - j_0^2 e^{-2u}) u_x (u - u_D)_x dx \\
 & = \int_0^1 \mathcal{A}^{-2} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) (u_x^2 - u_x u_{D,x}) dx \\
 & \quad + \int_0^1 (p''(\mathcal{A}) \mathcal{A} + 2j_0^2 \mathcal{A}^{-2}) (u - u_D) u_x (u - u_D)_x dx \\
 & \quad + \frac{1}{2} \int_0^1 (p'''(e^y) e^{2y} + p''(e^y) e^y - 4j_0^2 e^{-2y}) u_x (u - u_D)_x (u - u_D)^2 dx \\
 & \geq \int_0^1 \mathcal{A}^{-2} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) (u_x^2 - u_x u_{D,x}) dx \\
 & \quad - \frac{9}{8} \max_{x \in [0,1]} (|p''(\mathcal{A})| \mathcal{A} + 2j_0^2 \mathcal{A}^{-2}) \delta_1 \int_0^1 u_x^2 dx \\
 & \quad - \frac{9}{16} \max_{\ln \frac{16}{25} \mathcal{A}_* \leq y \leq \ln \frac{25}{16} \mathcal{A}^*} (|p'''(e^y)| e^{2y} + |p''(e^y) e^y| + 4j_0^2 e^{-2y}) \delta_1^2 \int_0^1 u_x^2 dx \\
 & \quad - K_c \delta_0,
 \end{aligned}$$

where  $y = u + \theta_1(u - u_D)$  for some  $\theta_1 \in (0, 1)$ . This implies, using (3.7)–(3.9) and (3.12),

$$\begin{aligned}
 & \int_0^1 (p'(e^u) - j_0^2 e^{-2u}) u_x (u - u_D)_x dx \\
 & \geq \int_0^1 \mathcal{A}^{-2} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) (u_x^2 - u_x u_{D,x}) dx - \frac{1}{2} (1 + \theta) A_0 \\
 (3.14) \quad & \times \int_0^1 u_{xx}^2 dx - K_c \delta_0.
 \end{aligned}$$

Furthermore, it follows from (1.16)

$$\begin{aligned}
 & \int_0^1 (p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2}) (u_x^2 - u_x u_{D,x}) dx \\
 & \geq (1 + \theta) \min_{x \in E} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} \int_0^1 u_x^2 dx + (1 - \theta) \\
 & \quad \times \int_{I \setminus E} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} u_x^2 dx \\
 & \quad - \frac{1}{\theta} \max_{x \in [0,1]} |p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2}| \int_0^1 |u_{D,x}|^2 dx \\
 & \geq (1 + \theta) \min_{x \in [0,1]} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} \int_0^1 u_{xx}^2 dx + (1 - \theta) \\
 (3.15) \quad & \times \int_{I \setminus E} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} u_x^2 dx - \frac{1 + \kappa}{1 - \kappa} K_c \delta_0,
 \end{aligned}$$

where we have used

$$\int_E u_x^2 dx \leq \int_0^1 u_{xx}^2 dx.$$

The estimates (3.15) and (3.14) yield

$$\begin{aligned}
 & \int_0^1 (p'(e^u) - j_0^2 e^{-2u}) u_x (u - u_D)_x dx \\
 & \geq (1 + \theta) \min_{x \in [0,1]} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} \int_0^1 u_{xx}^2 dx - \frac{1}{2} (1 + \theta) A_0 \int_0^1 u_{xx}^2 dx \\
 (3.16) \quad & + (1 - \theta) \int_{I \setminus E} \left( p'(\mathcal{A}) - \frac{j_0^2}{\mathcal{A}^2} \right) u_x^2 dx - \frac{1 + \kappa}{1 - \kappa} K_c \delta_0.
 \end{aligned}$$

Substituting (3.13) and (3.16) into (3.11) and using (3.10), we have

$$\begin{aligned}
 & A_0 \|u_{xx}\|^2 + \mathcal{A}_* \|u - u_D\|^2 \\
 (3.17) \quad & + \int_{I \setminus E} (p'(\mathcal{A}) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} (u - u_D)_x^2 dx \leq \frac{1 + \kappa}{1 - \kappa} K_c \delta_0,
 \end{aligned}$$

where we recall that  $\mathcal{A}_*$  and  $A_0$  are given by (3.3) and (3.4), respectively.



Now, we turn to higher order estimates. Let  $u \in H^4$  be a solution to the BVP (1.12). Multiply this equation with  $\varepsilon^2 u_{xxxx}$  and integrate over  $(0, 1)$  to obtain

$$\begin{aligned}
 & \frac{1}{4} \varepsilon^4 \int_0^1 [u_{xxxx}^2 - (u_x u_{xxx} + u_{xx}^2) u_{xxxx}] dx \\
 &= \varepsilon^2 \int_0^1 ((p'(e^u) - j_0^2 e^{2u}) u_x)_x u_{xxxx} dx - \varepsilon^2 \int_0^1 (e^u - e^{u_D}) u_{xxxx} dx \\
 (3.18) \quad & - \varepsilon^2 \int_0^1 (\mathcal{A} - \mathcal{C}) u_{xxxx} dx - \varepsilon^2 \frac{j_0}{\tau} \int_0^1 e^{-u} u_x u_{xxxx} dx.
 \end{aligned}$$

Due to (1.13), there are  $y_1, y_2, y_3, y_4 \in (0, 1)$  such that

$$u_x(y_1) = u_{xx}(y_2) = u_{xx}(y_3) = u_{xxx}(y_4) = 0,$$

and

$$(3.19) \quad u_{xx}^2(x) + \int_0^1 u_{xx}^2 dx \leq \int_0^1 u_{xxx}^2 dx, \quad \int_0^1 u_{xxx}^2 dx \leq \int_0^1 u_{xxxx}^2 dx.$$

Thus, it follows from (3.17), (3.19) and Hölder's inequality

$$\begin{aligned}
 & \int_0^1 (u_x u_{xxx} + u_{xx}^2) u_{xxxx} dx \\
 & \leq |u_x|_\infty \|u_{xxx}\| \cdot \|u_{xxxx}\| + |u_{xx}|_\infty \|u_{xx}\| \cdot \|u_{xxxx}\| \\
 (3.20) \quad & \leq \frac{1 + \kappa}{1 - \kappa} K_c \delta_0 \int_0^1 u_{xxxx}^2 dx.
 \end{aligned}$$

Then, we obtain from (3.18), in view of (3.7), (3.17), (3.20), (3.19), and Cauchy's inequality, that

$$\begin{aligned}
 \varepsilon^4 \int_0^1 u_{xxxx}^2 dx & \leq K_c \int_0^1 [(p'(e^u) - j_0^2 e^{-2u})^2 u_{xx}^2 + (p''(e^u) e^u + 2j_0^2 e^{-2u})^2 u_x^4] dx \\
 (3.21) \quad & + K_c \int_0^1 [(u - u_D)^2 + (\mathcal{A} - \mathcal{C})^2 + u_x^2] dx \leq \frac{1 + \kappa}{1 - \kappa} K_c \delta_0,
 \end{aligned}$$

provided that  $\delta_0$  is small enough. By (3.19) and (3.21), we have

$$(3.22) \quad \varepsilon^4 \int_0^1 [u_{xxx}^2 + u_{xxxx}^2] dx \leq \frac{1 + \kappa}{1 - \kappa} K_c \delta_0.$$

The combination of (3.17) and (3.22) finally leads to

$$\begin{aligned}
 & \mathcal{A}_* \|u - u_D\|^2 + A_0 \|u_{xx}\|^2 + \varepsilon^4 \|(u_{xxx}, u_{xxxx})\|^2 \\
 (3.23) \quad & + \int_{\Gamma \setminus E} (p'(\mathcal{A}^2) \mathcal{A}^2 - j_0^2) \mathcal{A}^{-2} (u - u_D)_x^2 dx \leq \frac{1 + \kappa}{1 - \kappa} K_c \delta_0.
 \end{aligned}$$

*Step 2. Existence.* It is not difficult to prove that there exists a solution  $u \in H^4$  to the BVP (1.12)–(1.13). The argument is similar to that used in section 2

based on the Leray-Schauder fixed point theorem. The function space is  $X := C^{0,1}([0, 1])$ . The corresponding linear BVP is

$$\begin{aligned} & \frac{1}{4}\varepsilon^2(u_{xx} + \frac{1}{2}\lambda v_x^2)_{xx} - (p'(e^{v_K})u_x - \lambda j_0^2 e^{-2v_K} v_x)_x \\ & \quad + \lambda\left(\frac{e^v - 1}{v}u + 1 - \mathcal{C}\right) = \lambda \frac{j_0}{\tau}(e^{-v_K})_x, \\ & u(0) = \lambda u_1, \quad u(1) = \lambda u_2, \quad u_x(0) = u_x(1) = 0, \end{aligned}$$

where  $\lambda \in [0, 1]$ ,  $v \in X$  and  $v_K = \min\{\delta_1 \ln \mathcal{A}^*, \max\{-\delta_1 \ln \mathcal{A}_*, v\}\}$  with  $\delta_1$  chosen such that (3.7)–(3.9) hold. The bilinear form and functional are defined respectively by

$$a(u, \psi) = \int_0^1 \left( \frac{1}{4}\varepsilon^2 u_{xx} v_{xx} + p'(e^{v_K})u_x \psi_x + \lambda \frac{e^v - 1}{v} u \psi \right) dx$$

and

$$F(\psi) = \int_0^1 \left( -\frac{1}{8}\varepsilon^2 \lambda v_x^2 \psi_{xx} - \lambda j_0^2 e^{-2v_K} v_x \psi_x + \lambda(\mathcal{C} - 1)\psi - \lambda \frac{j_0}{\tau} e^{-v_K} \psi \right) dx.$$

where  $v_K$ . We omit the details.

*Step 3. Uniqueness.* Let  $u, v \in H^4$  be two solutions to the BVP (1.12)–(1.13) satisfying (3.23). Using  $u - v \in H_0^2$  as an admissible test function in the weak formulation derived for  $u - v$ , we have

$$\begin{aligned} & \frac{1}{4}\varepsilon^2 \int_0^1 (u - v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x + v_x)(u - v)_x (u - v)_{xx} dx \\ & \quad + \int_0^1 ((p'(e^u) - j_0^2 e^{-2u})u_x - (p'(e^v) - j_0^2 e^{-2v})v_x) (u - v)_x dx \\ (3.24) \quad & = - \int_0^1 (e^u - e^v)(u - v) dx + \frac{j_0}{\tau} \int_0^1 (e^{-u} - e^{-v})_x (u - v) dx. \end{aligned}$$

The last term in (3.24) can be estimated as

$$\begin{aligned} & \frac{j_0}{\tau} \int_0^1 (e^{-u} - e^{-v})_x (u - v) dx \\ & = - \frac{j_0}{\tau} \int_0^1 [(e^{-u} - e^{-v})(u - v)u_x + \frac{1}{2}e^{-v}[(u - v)^2]_x] dx \\ (3.25) \quad & \leq \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0 \int_0^1 (u - v)^2 dx, \end{aligned}$$

where we have used

$$(3.26) \quad |(u_x, v_x)|_\infty \leq \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0,$$

derived from (3.23), and  $K_c > 0$  is a generic constant depending on  $\mathcal{A}$ ,  $\tau$ , and  $j_0$ . From (3.7), (3.24), and (3.25) follows

$$\begin{aligned}
 & \frac{1}{4}\varepsilon^2 \int_0^1 (u-v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x+v_x)(u-v)_x(u-v)_{xx} dx \\
 & \quad + \int_0^1 ((p'(e^u) - j_0^2 e^{-2u})u_x - (p'(e^v) - j_0^2 e^{-2v})v_x)(u-v)_x dx \\
 & \leq - \left( \frac{3}{4}\mathcal{A}_* - \sqrt{\frac{1+\kappa}{1-\kappa} \frac{K_c}{A_0} \delta_0} \right) \int_0^1 (u-v)^2 dx \\
 (3.27) \quad & \leq -\frac{1}{2}\mathcal{A}_* \int_0^1 (u-v)^2 dx,
 \end{aligned}$$

provided that  $\delta_0$  is small enough.

By (3.7), (3.26), Hölder’s inequality and

$$(3.28) \quad \int_0^1 (u-v)_x^2 dx \leq \int_0^1 (u-v)_{xx}^2 dx,$$

we have

$$\begin{aligned}
 & \frac{1}{4}\varepsilon^2 \int_0^1 (u-v)_{xx}^2 dx + \frac{1}{8}\varepsilon^2 \int_0^1 (u_x+v_x)(u-v)_x(u-v)_{xx} dx \\
 (3.29) \quad & \geq \frac{1}{4}\varepsilon^2 \left( 1 - \sqrt{\frac{1+\kappa}{1-\kappa} \frac{K_c}{A_0} \delta_0} \right) \int_0^1 (u-v)_{xx}^2 dx.
 \end{aligned}$$

By (1.16), (3.7) and (3.28), the third term on the left-hand side of (3.27) can be estimated as follows, using an approach similar to (3.14):

$$\begin{aligned}
 & \int_0^1 ((p'(e^u) - j_0^2 e^{-2u})u_x - (p'(e^v) - j_0^2 e^{-2v})v_x)(u-v)_x dx \\
 & = \int_0^1 (p'(e^u) - j_0^2 e^{-2u})(u-v)_x^2 dx \\
 & \quad + \int_0^1 (p'(e^u) - p'(e^v) - j_0^2 e^{-2u} + j_0^2 e^{-2v})v_x(u-v)_x dx \\
 & \geq \min_{x \in E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2} \int_0^1 (u-v)_x^2 dx \\
 & \quad + \int_{I \setminus E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2}(u-v)_x^2 dx \\
 & \quad + \int_0^1 (p'(e^u) - j_0^2 e^{-2u} - p'(\mathcal{A}) + j_0^2 \mathcal{A}^{-2})(u-v)_x^2 dx \\
 & \quad - |v_x|_\infty \int_0^1 |p'(e^u) - p'(e^v) - j_0^2 e^{-2u} + j_0^2 e^{-2v}|| (u-v)_x| dx \\
 & \geq \min_{x \in E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2} \int_0^1 (u-v)_{xx}^2 dx
 \end{aligned}$$

$$\begin{aligned}
 & + \int_{I \setminus E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2}(u - v)_x^2 dx \\
 & - \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0 \int_0^1 [(u - v)^2 + (u - v)_{xx}^2] dx \\
 & \geq \min_{x \in E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2} \int_0^1 (u - v)_{xx}^2 dx \\
 & + \int_{I \setminus E} (p'(\mathcal{A})\mathcal{A}^2 - j_0^2)\mathcal{A}^{-2}(u - v)_x^2 dx \\
 (3.30) \quad & - \frac{1}{3} \mathcal{A}_* \int_0^1 (u - v)^2 dx - \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0 \int_0^1 (u - v)_{xx}^2 dx,
 \end{aligned}$$

provided that  $\delta_0$  is small enough.

Substituting (3.29) and (3.30) into (3.27) leads to

$$\begin{aligned}
 & \min_{x \in [0,1]} \left( \frac{1}{4} \kappa \varepsilon^2 + p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2} \right) \int_0^1 (u - v)_{xx}^2 dx \\
 & + \left( \frac{1}{4} \varepsilon^2 (1 - \kappa) - \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0 \right) \int_0^1 (u - v)_{xx}^2 dx \\
 (3.31) \quad & + \frac{1}{6} \mathcal{A}_* \int_0^1 (u - v)^2 dx \leq 0,
 \end{aligned}$$

which implies that  $u = v$  in  $(0, 1)$  if  $\delta_0$  is so small that

$$\frac{1}{4} \varepsilon^2 (1 - \kappa) - \sqrt{\frac{1 + \kappa K_c}{1 - \kappa A_0}} \delta_0 > 0$$

and if the condition

$$(3.32) \quad \min_{x \in [0,1]} \left\{ \frac{1}{4} \kappa \varepsilon^2 + p'(\mathcal{A}) - j_0^2 \mathcal{A}^{-2} \right\} > 0$$

holds. The proof of Theorem 3.1 is completed. □

The existence and uniqueness of stationary solutions of (1.7)–(1.11) follows immediately from Theorem 3.1:

**Theorem 3.3.** *Assume that (1.15) and (1.16) hold and that  $p \in C^3(0, \infty)$ . For  $\kappa \in (0, 1)$ , assume that it holds*

$$(3.33) \quad \min_{x \in [0,1]} \mathcal{A}(x)^2 \left( \frac{1}{4} \kappa \varepsilon^2 + p'(\mathcal{A}(x)) \right) > j_0^2.$$

Then, the BVP (1.7)–(1.11) has a unique solution  $(\rho_0, \phi_0) \in H^4 \times H^2$  such that

$$(3.34) \quad \mathcal{A}_* \|\rho_0 - \mathcal{A}\|^2 + A_0 \|\rho_{0xx}\|^2 + \varepsilon^4 \|(\rho_{0xxx}, \rho_{0xxxx})\|^2 + \|\phi_{0x}\|_1^2 \leq \tilde{K}_c \delta_0,$$

provided that  $\|\mathcal{A}'\|_1 + \|\mathcal{A} - C\|$  is sufficiently small. The constant  $A_0$  is defined in (3.4) and  $\tilde{K}_c > 0$  is a constant depending on  $j_0, \tau$  and  $\mathcal{A}$ .

**Remark 3.4.** The estimates of Theorem 2.1 and Theorem 3.1 show that one can pass to the limit  $\varepsilon \rightarrow 0$  in the quantum Euler-Poisson system (similarly as in [16]) to obtain a solution to the classical Euler-Poisson system:

$$(3.35) \quad \partial_t \rho + j_x = 0,$$

$$(3.36) \quad \partial_t j + \left( \frac{j^2}{\rho} + p(\rho) \right)_x - \rho \phi_x - \frac{\varepsilon^2}{2} \rho \left( \frac{(\sqrt{\rho})_{xx}}{\sqrt{\rho}} \right)_x = -\frac{j}{\tau},$$

$$(3.37) \quad \lambda^2 \phi_{xx} = \rho - \mathcal{C}(x) \quad \text{in } [0, 1] \times (0, T),$$

The solution of this system (together with appropriate boundary conditions) is classical because the condition (2.1) or (3.1) reduces to the classical subsonic condition for (3.35)–(3.37). For the mathematical analysis on the Euler-Poisson system, we refer to [2, 17, 24] and references therein.

The following theorem is important for the stability analysis of stationary solutions obtained by Theorem 3.3 (see [22] for details).

**Theorem 3.5.** *Let  $(\rho_0, j_0, \phi_0)$  the unique strong solution given by Theorem 3.3. Let  $\omega_0 = \sqrt{\rho_0}$ . Then  $(\omega_0, j_0, \phi_0)$  is the unique solution of the following BVP*

$$(3.38) \quad j_x = 0,$$

$$(3.39) \quad \left( \frac{j^2}{w^2} + p(w^2) \right)_x = w^2 \phi_x + \frac{1}{2} \varepsilon^2 w^2 \left( \frac{w_{xx}}{w} \right)_x - \frac{j}{\tau},$$

$$(3.40) \quad \phi_{xx} = w^2 - \mathcal{C}(x),$$

with boundary conditions

$$(3.41) \quad w(0) = \sqrt{\rho_1}, \quad w(1) = \sqrt{\rho_2}, \quad w_x(0) = w_x(1) = 0,$$

$$(3.42) \quad \phi(0) = 0, \quad \phi(1) = \Phi_0.$$

Moreover, it holds

$$(3.43) \quad \|\omega_0 - \sqrt{\mathcal{A}}\|^2 + \|\omega_{0x}\|_3^2 + \|\phi_{0x}\|_1^2 \leq \tilde{C}_0 \delta_0,$$

where  $\tilde{C}_2 > 0$  is a constant depending on  $\mathcal{A}$ ,  $j_0$ ,  $\tau$  and  $\varepsilon$ .

**Acknowledgements.** The first author acknowledges partial support from the Deutsche Forschungsgemeinschaft, grants JU 359/3 (Gerhard-Hess Award) and JU 359/5 (Priority Program “Multi-Scale Problems”), and by the EU Project “Hyperbolic and Kinetic Equations”, grant HPRN-CT-2002-00282. The second acknowledges the partial support from the Wittgenstein Award 2000 of P. Markowich funded by the Austrian FWF and from the Austrian-Chinese Scientific-Technical Collaboration Agreement.

## REFERENCES

- [1] Brezzi, F., Gasser, I., Markowich, P., Schmeiser, C., *Thermal equilibrium state of the quantum hydrodynamic model for semiconductor in one dimension*, Appl. Math. Lett. **8** (1995), 47–52.
- [2] Chen, G. and Wang, D., *Convergence of shock schemes for the compressible Euler-Poisson equations*, Comm. Math. Phys. **179** (1996), 333–364.
- [3] Courant, R. and Friedrichs, K. O., *Supersonic flow and shock waves*, Springer-Verlag, New York 1976.
- [4] Degond, P. and Markowich, P. A., *On a one-dimensional steady-state hydrodynamic model*, Appl. Math. Lett. **3** (1990), 25–29.
- [5] Degond, P. and Markowich, P. A., *A steady state potential flow model for semiconductors*, Ann. Mat. Pura Appl. **165** (1993), 87–98.
- [6] Gamba, I., *Stationary transonic solutions of a one-dimensional hydrodynamic model for semiconductor*, Comm. Partial Differential Equations **17** (1992), 553–577.
- [7] Gamba, I. and Jüngel, A., *Asymptotic limits in quantum trajectory models*, Comm. Partial Differential Equations **27** (2002), 669–691.
- [8] Gamba, I. and Jüngel, A., *Positive solutions to singular second and third order differential equations for quantum fluids*, Arch. Rational Mech. Anal. **156** (2001), 183–203.
- [9] Gamba, I. and Morawitz, C., *A viscous approximation for a 2D steady semiconductor or transonic gas dynamics flow: existence theorem for potential flow*, Comm. Pure Appl. Math. **49** (1996), 999–1049.
- [10] Gardner, C., *Numerical simulation of a steady-state electron shock wave in a submicron semiconductor device*, IEEE Trans. El. Dev. **38** (1991), 392–398.
- [11] Gardner, C., *The quantum hydrodynamic model for semiconductors devices*, SIAM J. Appl. Math. **54** (1994), 409–427.
- [12] Gasser, I. and Jüngel, A., *The quantum hydrodynamic model for semiconductors in thermal equilibrium*, Z. Angew. Math. Phys. **48** (1997), 45–59.
- [13] Gasser, I., Lin, C.-K. and Markowich, P., *A review of dispersive limits of the (non)linear Schrödinger-type equation*, Taiwanese J. of Math. **4**, (2000), 501–529.
- [14] Gasser, I. and Markowich, P., *Quantum hydrodynamics, Wigner transforms and the classical limit*, Asymptot. Anal. **14** (1997), 97–116.
- [15] Gasser, I., Markowich, P. A. and Ringhofer, C., *Closure conditions for classical and quantum moment hierarchies in the small temperature limit*, Transport Theory Statistic Phys. **25** (1996), 409–423.
- [16] Gyi, M. T. and Jüngel, A., *A quantum regularization of the one-dimensional hydrodynamic model for semiconductors*, Adv. Differential Equations **5** (2000), 773–800.
- [17] Hsiao, L. and Yang, T., *Asymptotic of initial boundary value problems for hydrodynamic and drift diffusion models for semiconductors*, J. Differential Equations **170** (2001), 472–493.
- [18] Jerome, J., *Analysis of charge transport: a mathematical study of semiconductor devices*, Springer-Verlag, Heidelberg 1996.
- [19] Jüngel, A., *A steady-state potential flow Euler-Poisson system for charged quantum fluids*, Comm. Math. Phys. **194** (1998), 463–479.
- [20] Jüngel, A., *Quasi-hydrodynamic semiconductor equations*, Progress in Nonlinear Differential Equations, Birkhäuser, Basel 2001.
- [21] Jüngel, A., Mariani, M. C. and Rial, D., *Local existence of solutions to the transient quantum hydrodynamic equations*, Math. Models Methods Appl. Sci. **12** (2002), 485–495.
- [22] Jüngel, A. and Li, H.-L., *Quantum Euler-Poisson systems: global existence and exponential decay*, to appear in Quart. Appl. Math. 2005.

- [23] Landau, L. D. and Lifshitz, E. M., *Quantum mechanics: non-relativistic theory*, New York, Pergamon Press 1977.
- [24] Li, H.-L. and Markowich, P. A., *A review of hydrodynamical models for semiconductors: asymptotic behavior*, Bol. Soc. Brasil. Mat. (N.S.) **32** (2001), 321-342.
- [25] Loffredo, M. and Morato, L., *On the creation of quantum vortex lines in rotating HeII*. Il nuovo cimento **108B** (1993), 205-215.
- [26] Madelung, E., *Quantentheorie in hydrodynamischer Form*, Z. Phys. **40** (1927), 322.
- [27] Marcati, P. and Natalini, R., *Weak solutions to a hydrodynamic model for semiconductors and relaxation to the drift-diffusion equation*, Arch. Rational Mech. Anal. **129** (1995), 129-145.
- [28] Markowich, P., Ringhofer, C. and Schmeiser, C., *Semiconductor Equations*, Springer, Wien 1990.
- [29] Pacard, F. and Unterreiter, A., *A variational analysis of the thermal equilibrium state of charged quantum fluids*, Comm. Partial Differential Equations **20** (1995), 885-900.
- [30] Shu, C.-W., *Essentially non-oscillatory and weighted essentially non-oscillatory schemes for hyperbolic conservation laws*, ICASE Report No. 97-65, NASA Langley Research Center, Hampton, USA 1997.
- [31] Zhang, B. and Jerome, W., *On a steady-state quantum hydrodynamic model for semiconductors*, Nonlinear Anal., Theory Methods Appl. **26** (1996), 845-856.

†FACHBEREICH MATHEMATIK UND INFORMATIK, UNIVERSITÄT MAINZ  
STAUDINGERWEG 9, 55099 MAINZ, GERMANY  
*E-mail:* juengel@mathematik.uni-mainz.de

\*DEPARTMENT OF MATHEMATICS, CAPITAL NORMAL UNIVERSITY  
BEIJING 100037, P. R. CHINA  
*E-mail:* lih1@math.sci.osaka-u.ac.jp