

# A Model of Heat Conduction

P. Collet<sup>1</sup>, J.-P. Eckmann<sup>2</sup>

<sup>1</sup> Centre de Physique Théorique, CNRS UMR 7644, Ecole Polytechnique,  
F-91128 Palaiseau Cedex, France

<sup>2</sup> Département de Physique Théorique et Section de Mathématiques,  
Université de Genève, CH-1211 Genève 4, Switzerland.  
E-mail: Jean-Pierre.Eckmann@physics.unige.ch

Received: 18 April 2008 / Accepted: 8 August 2008  
Published online: 21 November 2008 – © Springer-Verlag 2008

**Abstract:** In this paper, we first define a deterministic particle model for heat conduction. It consists of a chain of  $N$  identical subsystems, each of which contains a scatterer and with particles moving among these scatterers. Based on this model, we then derive heuristically, in the limit of  $N \rightarrow \infty$  and decreasing scattering cross-section, a Boltzmann equation for this limiting system. This derivation is obtained by a closure argument based on memory loss between collisions. We then prove that the Boltzmann equation has, for stochastic driving forces at the boundary, close to Maxwellians, a unique non-equilibrium steady state.

## 1. Introduction

In this paper, we consider the problem of heat conduction for the continuum limit ( $N \rightarrow \infty$ ) of a particle model of a chain of  $N$  cells, each of which contains a (very simple) scatterer in its interior. Particles move between the cells, interacting with the scatterers, but not among themselves, similar to the model put forward in [4].

The paper begins with a detailed description of the scatterer model, which is one-dimensional, with the dynamics also in 1 dimension. Once the principles of interaction are laid down, we proceed to derive, in a heuristic way, under the usual closure assumptions, a Boltzmann equation for the limit system. (The reader who is interested only in the formulation of the Boltzmann equation can directly skip to Eq. (4.5).) It should be noted that in order to obtain a reasonable formal limit of the particle system, the scattering probabilities have to be chosen proportional to  $1/N$ .

Our main object of mathematical study is then this Boltzmann equation. In particular, starting from Sec. 5 we show how to formulate the heat conduction problem. Namely, heat conduction problems are usually described by prescribing the incoming fluxes on both ends of the system. We however reformulate this as a problem of prescribing what happens at one end only and then to find out what happens at the other. Using the inverse function theorem on a suitable space, we will be able to show that the problem is related

to a diffeomorphism in a suitable space (in fact a Banach cone) (starting from Sec. 7). This Banach cone allows for a wide range of initial conditions, and is basically only limited by conditions on the velocity distributions at infinite momenta. However, pure Maxwellians are not in this cone (but on its boundary), although they naturally form the equilibrium solutions. In a forthcoming paper with C. Mejía-Monasterio [1], we will illustrate some of the properties of both the particle model and its Boltzmann formulation, and also “justify” numerically the approximations which lead to the Boltzmann equation.

## 2. The Particle Model

**2.1. One cell.** To define the particle model, we begin by describing the scattering process in one cell. We begin with the description of one “cell”. The cell is 1-dimensional, of length  $2L$ , and with particles entering on either side. These particles all have the same mass  $m$ , some velocity  $v$  and momentum  $p = mv$ . These particles do not interact among themselves. Note that  $v \in \mathbb{R}$  and more precisely,  $v > 0$  if the particle enters from the left, while  $v < 0$  if it enters on the right side of the cell. In the center of the cell, we imagine a “scatterer” which is a point-like particle which can exchange energy and momentum with the particles, but does not change its own position. (This scatterer is to be thought of as a 1-dimensional variant of the rotating disks used in [4].) The scatterer has mass  $M$  and its “velocity” will be denoted by  $V$ . The collision rules are those of an elastic collision, where  $\tilde{v}$  and  $\tilde{V}$  denote quantities after the collision while  $v, V$  are those before the collision. In equations,

$$\begin{aligned} \tilde{v} &= -\varrho v + (1 + \varrho)V, \\ \tilde{V} &= (1 - \varrho)v + \varrho V, \end{aligned}$$

with

$$\varrho = \frac{M - m}{M + m}, \quad \mu \equiv \frac{m}{M} = \frac{1 - \varrho}{1 + \varrho}. \tag{2.1}$$

Note that  $\varrho \in (-1, 1)$ , since we assume  $m$  and  $M$  to be finite and non-zero. If  $\tilde{v} > 0$ , we say that the particle leaves the cell to the right; if  $\tilde{v} < 0$ , we say it leaves to the left. In particular, although everything is 1-dimensional, particles *can* cross the scatterers.

For simplicity, we will assume  $\varrho > 0$ , that is,  $M > m$ . For the momenta, we get the analogous rules

$$\begin{aligned} \tilde{p} &= -\varrho p + (1 - \varrho)P, \\ \tilde{P} &= (1 + \varrho)p + \varrho P. \end{aligned}$$

Note that the matrix

$$S = \begin{pmatrix} -\varrho & 1 - \varrho \\ 1 + \varrho & \varrho \end{pmatrix} \tag{2.2}$$

has determinant equal to  $-1$  and furthermore  $S^2 = 1$ .

We next formulate scattering in terms of probability densities (for momenta) for just one cell. We denote by  $g(t, P)$  the probability density that at time  $t$  the scatterer has momentum  $P$  ( $= MV$ ) and we will establish the equation for the time evolution of this function. To begin with, we assume that particles enter only from the left of the cell, with

momentum distribution (in a neighboring cell or a bath)  $p \mapsto f_L^+(t, p)$ , where  $p = mv$ . Thus, there are, on average,  $p f_L^+(t, p) dp/m$  particles entering the cell (per unit of time) from the left with momentum in  $[p, p + dp]$ .<sup>1</sup> Note that  $f_L^+$  has support on  $p \geq 0$  only, (indicated by the exponent “+”); it is the distribution of particles *going to enter the cell* from the left. Also note that the distribution of the momenta after collision, *i.e.*, before *leaving* the cell, is in general not the same as  $f_L^+$ .

Denote by  $\mathcal{P}(t)$  the stochastic process describing the momentum of the scatterer. We have for any interval (measurable set) of momenta  $A$ , for the probabilities  $\mathbf{P}$ :

$$\begin{aligned} \mathbf{P}(\mathcal{P}(t + dt) \in A) &= \mathbf{P}(\mathcal{P}(t) \in A; \text{ no collision in } [t, t + dt]) \\ &\quad + \mathbf{P}(\mathcal{P}(t + dt) \in A; \text{ collisions occurred in } [t, t + dt]). \end{aligned}$$

We assume for simplicity that with probability one, only one collision can occur in an interval  $[t, t + dt]$ . If there is a collision in  $[t, t + dt]$  with a particle of velocity  $v = p/m > 0$ , this particle must have left the boundary at time  $t - \frac{m}{p}L$  with momentum  $p$ . Therefore,

$$\begin{aligned} &\mathbf{P}(\mathcal{P}(t) \in A; \text{ a collision occurred in } [t, t + dt]) \\ &= dt \int_A d\tilde{P} \int_{\mathbb{R}^+} dp \delta\left(\tilde{P} - \varrho P - (1 + \varrho)p\right) g(t, P) \frac{p}{m} f_L^+\left(t - \frac{m}{p}L, p\right). \end{aligned}$$

This formulation neglects memory effects coming from the fact that a particle may have hit the scatterer, bounce out of the cell and reenter to hit again the scatterer. Similarly,

$$\begin{aligned} &\mathbf{P}(\mathcal{P}(t) \in A; \text{ no collision occurred in } [t, t + dt]) \\ &= \left(1 - dt \int_{\mathbb{R}^+} dp \frac{p}{m} f_L^+\left(t - \frac{m}{p}L, p\right)\right) \int_A dP g(t, P). \end{aligned}$$

We immediately deduce the evolution equation,

$$\begin{aligned} \partial_t g(t, P) &= -g(t, P) \int_{\mathbb{R}^+} dp \frac{p}{m} f_L^+\left(t - \frac{m}{p}L, p\right) \\ &\quad + \frac{1}{\varrho} \int_{\mathbb{R}^+} dp g\left(t, \frac{P - (1 + \varrho)p}{\varrho}\right) \frac{p}{m} f_L^+\left(t - \frac{m}{p}L, p\right). \end{aligned} \tag{2.3}$$

Note that this equation preserves the integral of  $g$  over  $P$ , *i.e.*, it preserves probability.

This identity generalizes immediately to the inclusion of injection from the right, with distribution  $f_R^-$  having support in  $p < 0$ . One gets

$$\begin{aligned} \partial_t g(t, P) &= -g(t, P) \int_{\mathbb{R}} dp \frac{|p|}{m} \left(f_L^+\left(t - \frac{m}{p}L, p\right) + f_R^-\left(t + \frac{m}{p}L, p\right)\right) \\ &\quad + \frac{1}{\varrho} \int_{\mathbb{R}} dp g\left(t, \frac{P - (1 + \varrho)p}{\varrho}\right) \frac{|p|}{m} \left(f_L^+\left(t - \frac{m}{p}L, p\right) + f_R^-\left(t + \frac{m}{p}L, p\right)\right). \end{aligned} \tag{2.4}$$

In the stationary case, this leads to

$$g(P) = \frac{1}{\varrho \lambda} \int_{\mathbb{R}} dp g\left(\frac{P - (1 + \varrho)p}{\varrho}\right) \frac{|p|}{m} \left(f_L^+(p) + f_R^-(p)\right), \tag{2.5}$$

---

<sup>1</sup> The factor  $p$  takes into account the probability of *crossing* the boundary of the cell (which is not the same as the probability of a particle with momentum  $p$  to *be* in the cell).

where

$$\lambda = \int_{\mathbb{R}} dp \frac{|p|}{m} (f_L^+(p) + f_R^-(p)) \tag{2.6}$$

is the particle flux (see Sec. 4.1 below).

It is important to note that the solution  $g$  of Eq. (2.5) only depends on the *sum*:  $f = f_L^+ + f_R^-$ , and thus, we can define a map

$$f \mapsto g_f,$$

where  $g_f$  is the (unique) solution of Eq. (2.5). It will be discussed in detail in Sec. 6.

We can also compute the distribution of the momenta of the particles after collision. We have

$$\begin{aligned} & \mathbf{P}(\tilde{p} \in A; \text{ a collision occurred in } [t, t + dt]) \\ &= dt \int_A d\tilde{p} \int_{\mathbb{R}} dP \delta(\tilde{p} + \varrho P - (1 - \varrho)P) g(t, P) \frac{|p|}{m} f\left(t - \frac{m}{|\tilde{p}|}L, p\right). \end{aligned}$$

This particle reaches the left or right boundary of the cell (according to the sign of  $\tilde{p}$ ) after a time  $mL/|\tilde{p}|$  (assuming the scatterer is located in the center of the cell). Therefore, we have for the ejection distributions  $f_L^-$  (on the left) and  $f_R^+$  (on the right):

$$\frac{|\tilde{p}|}{m} f_L^-(t, \tilde{p}) = \frac{\theta(-\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g\left(t - \frac{m}{|\tilde{p}|}L, \frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f\left(t - \frac{m}{|\tilde{p}|}L - \frac{m}{|\tilde{p}|}L, p\right),$$

and

$$\frac{|\tilde{p}|}{m} f_R^+(t, \tilde{p}) = \frac{\theta(+\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g\left(t - \frac{m}{|\tilde{p}|}L, \frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f\left(t - \frac{m}{|\tilde{p}|}L - \frac{m}{|\tilde{p}|}L, p\right),$$

where  $\theta$  is the Heaviside function. In the stationary case we get

$$\frac{|\tilde{p}|}{m} f_L^-(\tilde{p}) = \frac{\theta(-\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g\left(\frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f(p), \tag{2.7}$$

and

$$\frac{|\tilde{p}|}{m} f_R^+(\tilde{p}) = \frac{\theta(+\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g\left(\frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f(p). \tag{2.8}$$

Since  $g = g_f$  is determined by the incoming distribution  $f_{in} = f_L^+ + f_R^-$  (and is unique if we normalize the integral of  $g$  to 1)

$$\int_{\mathbb{R}} dP g(t, P) = 1, \tag{2.9}$$

we see that the outgoing distribution  $f_{out} = f_L^- + f_R^+$  is entirely determined by the incoming distribution. Note also that the flux is preserved:

$$\int dp \frac{|p|}{m} f_{in}(p) = \int dp \frac{|p|}{m} f_{out}(p).$$

2.2. *Stationary solutions for one cell.* Here, we will look for stationary states of the evolution equation (2.3), which have also the property that the ejected distributions are equal to the injected ones. It is almost obvious that Maxwellian fixed points can be found, but for completeness, we write down the formulas. The reader should note that the distributions  $f_{\text{in}}$  and  $f_{\text{out}}$  have singularities at  $p = 0$ . This reflects the well-known fact that slow particles need more time to leave the cell than fast ones. However  $F(p) \equiv \frac{p}{m} f(p)$  is a very nice function, and it is this function which appears in all the calculations of the fluxes, and stationary profiles. In this section we do the calculations with the quantity  $f$ . Starting from Sec. 4, we will use  $F$ .

We impose the two incoming distributions

$$f_{\text{L}}^+(p) = \sigma \theta(+p) \frac{m}{|p|} e^{-\beta p^2/(2m)},$$

and

$$f_{\text{R}}^-(p) = \sigma \theta(-p) \frac{m}{|p|} e^{-\beta p^2/(2m)},$$

where  $\sigma$  is an arbitrary positive constant (related to  $\lambda$  in (2.6)) and  $\theta$  is the Heaviside function. It is easy to verify, using Gaussian integration and the identity  $M = M\varrho^2 + m(1+\varrho)^2$ , that the solution of Eq. (2.5) is given by

$$g(P) = \sqrt{\frac{\beta}{2\pi M}} e^{-\beta P^2/(2M)} = \sqrt{\frac{\beta}{2\pi M}} e^{-\beta P^2(1-\varrho)/((1+\varrho)2m)}.$$

Moreover, using the same identity several times one gets from Eqs. (2.7) and (2.8) for the exiting distributions

$$f_{\text{L}}^-(p) = \sigma \theta(-p) \frac{m}{|p|} e^{-\beta p^2/(2m)},$$

and

$$f_{\text{R}}^+(p) = \sigma \theta(+p) \frac{m}{|p|} e^{-\beta p^2/(2m)}.$$

Therefore, we see that the Maxwellian fixed points (divided by  $|p|$ ) preserve both the distribution  $g$  of the scatterer, as well as the distributions of the particles.

In fact, there are also non-Maxwellian fixed points of the form

$$f_{\text{L}}^+(p) = \sigma \theta(+p) \frac{m}{|p|} e^{-\beta(p-ma)^2/(2m)},$$

and

$$f_{\text{R}}^-(p) = \sigma \theta(-p) \frac{m}{|p|} e^{-\beta(p-ma)^2/(2m)}.$$

It is easy to verify that the solution of Eq. (2.5) is now given by

$$g(P) = \sqrt{\frac{\beta}{2\pi M}} e^{-\beta(P-Ma)^2/(2M)}.$$

Moreover,

$$f_{\text{L}}^-(p) = \sigma \theta(-p) \frac{m}{|p|} e^{-\beta(p-ma)^2/(2m)},$$

and

$$f_{\text{R}}^+(p) = \sigma \theta(+p) \frac{m}{|p|} e^{-\beta(p-ma)^2/(2m)}.$$

The verification that this is a solution for any  $a \in \mathbb{R}$  is again by Gaussian integration. Note that if  $a \neq 0$  there is in fact a flux through the cell.

### 3. $N$ Cells

The model generalizes immediately to the case of  $N$  cells which are arranged in a row, by requiring that the exit distributions of any given cell are equal to the entry distributions of the neighboring cells: the cells are numbered from 1 to  $N$  and we have the collections of functions  $f_{L,i}^+$ ,  $f_{L,i}^-$ ,  $f_{R,i}^+$ , and  $f_{R,i}^-$  for the particle and  $g_i$  for the scatterers,  $i = 1, \dots, N$ . The equality of entrance and exit distributions is given by the identities  $f_{L,i+1}^+ = f_{R,i}^+$ , and  $f_{R,i}^- = f_{L,i+1}^-$  for  $1 \leq i < N$ . The system is completely determined by the two functions  $f_{L,1}^+$  and  $f_{R,N}^-$ . Equations (2.5) generalize to

$$g_i(P) = \frac{1}{\varrho\lambda} \int_{\mathbb{R}} dp g_i \left( \frac{P-(1+\varrho)p}{\varrho} \right) \frac{|p|}{m} \left( f_{L,i}^+(p) + f_{R,i}^-(p) \right), \tag{3.1}$$

and similarly (2.7) and (2.8) lead to

$$\begin{aligned} \frac{|\tilde{p}|}{m} f_{L,i}^-(\tilde{p}) &= \frac{\theta(-\tilde{p})}{1-\varrho} \int_{\mathbb{R}} dp g_i \left( \frac{\tilde{p}+\varrho p}{1-\varrho} \right) \frac{|p|}{m} f_i(p), \\ \frac{|\tilde{p}|}{m} f_{R,i}^+(\tilde{p}) &= \frac{\theta(+\tilde{p})}{1-\varrho} \int_{\mathbb{R}} dp g_i \left( \frac{\tilde{p}+\varrho p}{1-\varrho} \right) \frac{|p|}{m} f_i(p), \end{aligned} \tag{3.2}$$

where  $f_i = f_{L,i}^+ + f_{R,i}^-$ . Clearly, the Gaussians of the previous section are still solutions to the full equations for  $N$  contiguous cells.

Here we have closed the model by assuming independence between the particles leaving and entering from the left (and from the right). In concrete systems this is not true since a particle can leave a cell to the left and re-bounce back into the original cell after just one collision with the scatterer in the neighboring cell, and, in such a situation there is too much memory to allow for full independence. It is possible to imagine several experimental arrangements for which independence is a very good approximation, see also [3,4] for discussions of such issues. One of them could be to imagine long channels between the scatterers where time decorrelation would produce independence. Note that ‘‘chaotic’’ channels may be more complicated since they can modify the distribution of left (right) traveling particles between two cells.

### 4. Continuous Space

We are now ready to derive, non-rigorously, a continuum model. The cells are now replaced by a continuum, with a variable  $x \in [0, 1]$  replacing  $i/N$ , where  $i$  is the index of the  $i^{\text{th}}$  cell. The relations we have derived so far will be generalized to this continuum formulation. So we have moving particles, of mass  $m$  and described by a time-dependent density  $f(t, p, x)$ .

The scatterers have mass  $M$  and their momentum distribution is called  $g(t, P, x)$ . It is best to think that the continuous variable  $x \in [0, 1]$  replaces the discrete index  $i \in \{0, \dots, N\}$ . There is then an implicit rescaling of the form  $x \approx i/N$ . Recall that the scatterers are *fixed* in space (although they have momentum) but that the particles will move in the domain  $[0, 1]$ .

We first impose, for all  $x \in [0, 1]$ , the normalization

$$\int_{\mathbb{R}} dP g(t, P, x) = 1, \tag{4.1}$$

which is the generalization of Eq. (2.9). The particles again do not interact with each other, but only with the scatterers and, expressed in momenta, the matrix maps  $(p, P)$  to  $(\tilde{p}, \tilde{P})$ :

$$\begin{pmatrix} \tilde{p} \\ \tilde{P} \end{pmatrix} = \begin{pmatrix} -\varrho & 1 - \varrho \\ 1 + \varrho & \varrho \end{pmatrix} \begin{pmatrix} p \\ P \end{pmatrix} \equiv S \begin{pmatrix} p \\ P \end{pmatrix}, \tag{4.2}$$

as in Eq. (2.2).

We now rewrite the problem in the form of a Boltzmann equation, which takes into account this matrix, as well as the particle transport. One guesses, with  $(\tilde{p}, \tilde{P})$  related to  $(p, P)$  as above:

$$\begin{aligned} \partial_t f(t, p, x) + \frac{p}{m} \partial_x f(t, p, x) \\ = \int dP \left( \frac{|\tilde{p}|}{m} f(t, \tilde{p}, x) g(t, \tilde{P}, x) - \frac{|p|}{m} f(t, p, x) g(t, P, x) \right), \tag{4.3} \\ \partial_t g(t, P, x) = \int dp \left( \frac{|\tilde{p}|}{m} f(t, \tilde{p}, x) g(t, \tilde{P}, x) - \frac{|p|}{m} f(t, p, x) g(t, P, x) \right). \end{aligned}$$

The time independent version of the equation will be derived (non-rigorously) below from the model with a chain of cells. It is useful to introduce the function

$$F(t, p, x) = \frac{|p|}{m} f(t, p, x),$$

and then Eq. (4.3) takes the form

$$\begin{aligned} m \partial_t F(t, p, x) + p \partial_x F(t, p, x) \\ = |p| \int dP \left( F(t, \tilde{p}, x) g(t, \tilde{P}, x) - F(t, p, x) g(t, P, x) \right), \tag{4.4} \\ \partial_t g(t, P, x) = \int dp \left( F(t, \tilde{p}, x) g(t, \tilde{P}, x) - F(t, p, x) g(t, P, x) \right). \end{aligned}$$

*Remark.* One can also imitate a scattering cross section by introducing a factor  $\gamma \in [0, 1]$  in Eq. (4.4) (in the integrals) but this can be scaled away by a change of time and space scales. See also Sec. 9.3.

We come now to the **main equations** whose solutions will be discussed in detail in the remainder of the paper. Equation (4.4) takes, for the stationary solution, the form

$$\text{sign}(p) \partial_x F(p, x) = \int dP \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right), \tag{4.5a}$$

$$0 = \int dp \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right). \tag{4.5b}$$

We will show that this equation has non-equilibrium solutions.

*Remark.* Note that the model we have obtained here is *not* momentum-translation invariant, because of the term  $\text{sign}(p)$ , except when the r.h.s. of the equation is 0.

*Derivation of (4.5).* The derivation of (4.5) from (3.1–3.2) is based on the following formal limit: We replace the index  $i$  by the continuous variable  $x = i/N$  and set  $\varepsilon = 1/(2N)$ . We consider that  $f_{L,i}^\pm$  is at  $(i - \frac{1}{2})/N = x - \varepsilon$ , while  $f_{R,i}^\pm$  is at  $x + \varepsilon$ . We have the correspondences, with  $\theta_\pm(p) \equiv \theta(\pm p)$ :

$$\begin{aligned} \theta_+(p)F(p, x - \varepsilon) &= \frac{|p|}{m} f_{L,i}^+, & \theta_-(p)F(p, x - \varepsilon) &= \frac{|p|}{m} f_{L,i}^-, \\ \theta_+(p)F(p, x + \varepsilon) &= \frac{|p|}{m} f_{R,i}^+, & \theta_-(p)F(p, x + \varepsilon) &= \frac{|p|}{m} f_{R,i}^-, \\ g(P, x) &= g_i(P). \end{aligned}$$

To simplify momentarily the notation, let

$$\begin{aligned} F_-(p, x) \theta_+(p) &\equiv \theta_+(p)F(p, x - \varepsilon), & F_-(p, x) \theta_-(p) &\equiv \theta_-(p)F(p, x - \varepsilon), \\ F_+(p, x) \theta_+(p) &\equiv \theta_+(p)F(p, x + \varepsilon), & F_+(p, x) \theta_-(p) &\equiv \theta_-(p)F(p, x + \varepsilon). \end{aligned}$$

With these conventions, (3.1) becomes (setting  $\lambda = 1$ ):

$$g(P, x) = \frac{1}{\varrho} \int_{\mathbb{R}} dq g\left(\frac{P-(1+\varrho)q}{\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)), \tag{4.6}$$

which is equivalent to (4.5b). Similarly, Eq. (3.2) leads to

$$\begin{aligned} F_-(p, x)\theta_-(p) &= \frac{\theta(-p)}{1-\varrho} \int_{\mathbb{R}} dq g\left(\frac{p+\varrho q}{1-\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)), \\ F_+(p, x)\theta_+(p) &= \frac{\theta(+p)}{1-\varrho} \int_{\mathbb{R}} dq g\left(\frac{p+\varrho q}{1-\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)). \end{aligned} \tag{4.7}$$

Subtracting the first equation from the second in (4.7) leads to

$$\begin{aligned} F_+(p, x)\theta_+(p) - F_-(p, x)\theta_-(p) &= \frac{\text{sign}(p)}{1-\varrho} \int_{\mathbb{R}} dq g\left(\frac{p+\varrho q}{1-\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)). \end{aligned} \tag{4.8}$$

On the other hand, since  $dp dP = d\tilde{p} d\tilde{P}$  we can, by (4.5b), impose the condition

$$\int dP g(P, x) = 1,$$

for all  $x$ . Then we have the trivial identity

$$F_-(p, x)\theta_+(p) - F_+(p, x)\theta_-(p) = \int_{\mathbb{R}} dP g(P, x) (F_-(p, x)\theta_+(p) - F_+(p, x)\theta_-(p)). \tag{4.9}$$

Subtracting (4.9) from (4.8), we get for  $p > 0$ ,

$$\begin{aligned} F_+(p, x) - F_-(p, x) &= \frac{1}{1-\varrho} \int dq g\left(\frac{p+\varrho q}{1-\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)) \\ &\quad - \int dq g(q, x) F_-(p, x), \end{aligned} \tag{4.10}$$



while for  $p < 0$ ,

$$F_+(p, x) - F_-(p, x) = -\frac{1}{1-\varrho} \int dq g\left(\frac{p+\varrho q}{1-\varrho}, x\right) (F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q)) + \int dq g(q, x)F_+(p, x). \tag{4.11}$$

Note now that

$$\begin{aligned} &F_-(q, x)\theta_+(q) + F_+(q, x)\theta_-(q) \\ &= F(q, x - \varepsilon)\theta_+(q) + F(q, x + \varepsilon)\theta_-(q) \\ &= F(q, x) + \theta_+(q) (F(q, x - \varepsilon) - F(q, x)) \\ &\quad + \theta_-(q) (F(q, x + \varepsilon) - F(q, x)). \end{aligned}$$

Therefore, replacing the  $F_{\pm}$  in the r.h.s. in (4.10) and (4.11) by  $F(p, x)$  is a higher order correction in  $\varepsilon$ , and we finally get

$$F(p, x + \varepsilon) - F(p, x - \varepsilon) = \frac{\text{sign}(p)}{1-\varrho} \int_{\mathbb{R}} dq \left( g\left(\frac{p+\varrho q}{1-\varrho}, x\right) F(q, x) - g(q, x)F(p, x) \right).$$

A further change of integration variables leads to (4.5a), while (4.6) leads to (4.5b). (We have not taken into account another scaling of the scattering term by  $\varepsilon = 1/(2N)$  which is needed to get the derivative. This corresponds to a particle model where the scattering cross-section is proportional to  $1/N$ . (We will come back to this question in the discussion in Sect. 9.3 and in [1].) This ends the derivation of (4.5).

The derivative term in Eq. (4.5) reflects the gradients which have to appear when the system is out of equilibrium. However, if the system is at equilibrium, the equivalence between Eq. (4.5) and Eqs. (2.5)–(2.8) immediately tells us that stationary solutions in the form of Gaussians (for  $F$ , not for  $f$ ) exist:

$$F(p) = \gamma \sqrt{\frac{\beta}{2\pi m}} e^{-\beta p^2/(2m)}, \quad g(P) = \sqrt{\frac{\beta}{2\pi M}} e^{-\beta P^2/(2M)}. \tag{4.12}$$

Furthermore, we have again translated versions of this fixed point,

$$F(p) = \gamma \sqrt{\frac{\beta}{2\pi m}} e^{-\beta(p-ma)^2/(2m)}, \quad g(P) = \sqrt{\frac{\beta}{2\pi M}} e^{-\beta(P-Ma)^2/(2M)}, \tag{4.13}$$

because in this case, the r.h.s. of Eq. (4.5) is zero.

**4.1. Flux.** We can define various fluxes of the particles (recall that  $F(p, x) = |p| f(p, x)/m$ ):

$$\begin{aligned} \Phi_P &= \text{particle flux} &= \int dp \text{sign}(p) F(p, x), \\ \Phi_M &= \text{momentum activity} &= \int dp |p| F(p, x), \\ \Phi_E &= \text{energy flux} &= \int dp \frac{\text{sign}(p)p^2}{2m} F(p, x). \end{aligned} \tag{4.14}$$

Note that for the stationary Maxwellians of (4.13) these fluxes are equal to

$$\begin{aligned} \Phi_P &= \sqrt{\frac{2m\pi}{\beta}} \operatorname{erf}(a\sqrt{\beta m/2}), \\ \Phi_M &= \frac{2m}{\beta} e^{-\beta ma^2/2} + a\sqrt{\frac{2m^3\pi}{\beta}} \operatorname{erf}(a\sqrt{\beta m/2}), \\ \Phi_E &= \frac{2am^2}{\beta} e^{-\beta ma^2/2} + (1 + \beta ma^2)\sqrt{\frac{2m^3\pi}{\beta^3}} \operatorname{erf}(a\sqrt{\beta m/2}). \end{aligned}$$

Also note that for  $a = 0$  the quantity  $\Phi_M$  does not vanish. This is because it measures the total outgoing flux, not the directed outgoing flux (which is of course 0 when  $a = 0$ ).

**Lemma 4.1.** *For every stationary solution of (4.5) the 3 fluxes of (4.14) are independent of  $x \in [0, 1]$ .*

*Proof.* From (4.5a) we deduce that

$$\partial_x \int dp \operatorname{sign}(p) F(p, x) = \int dp dP \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right),$$

which vanishes since  $dp dP = d\tilde{p} d\tilde{P}$ . Similarly, multiplying (4.5a) by  $p$  and integrating over  $p$ , we get

$$\partial_x \int dp |p| F(p, x) = \int dp dP p \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right).$$

Multiplying (4.5b) by  $P$  and integrating over  $P$ , we get

$$0 = \int dp dP P \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right).$$

Adding these two equations, we see that

$$\partial_x \int dp |p| F(p, x) = \int dp dP (p + P) \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right).$$

But this vanishes, since  $P + p = \tilde{P} + \tilde{p}$  by momentum conservation, and using again  $dp dP = d\tilde{p} d\tilde{P}$ . In a similar way, we first have

$$\begin{aligned} \partial_x \int dp \frac{|p|p}{2m} F(p, x) \\ = \int dp dP \frac{p^2}{2m} \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right). \end{aligned}$$

Finally, multiplying this time (4.5b) by  $P^2/M$ , integrating over  $P$ , adding to the above equation and using energy conservation, we get

$$\begin{aligned} \partial_x \int dp \frac{|p|p}{2m} F(p, x) \\ = \int dp dP \left( \frac{p^2}{2m} + \frac{P^2}{2M} \right) \left( F(\tilde{p}, x) g(\tilde{P}, x) - F(p, x) g(P, x) \right) \\ = 0. \end{aligned}$$

Thus, all three fluxes are independent of  $x$ , as asserted.  $\square$

## 5. Formulating the Heat-Conduction Problem

Based on the stationarity equation (4.5), we now formulate the problem of heat conduction in mathematical terms. We imagine that the input of the problem is given by prescribing the *incoming* fluxes on both sides of the system. The system, in its stationary state, should then adapt all the other quantities,  $F$  and  $g$ , to this given input, which describes really the forcing of the system. In particular, the distribution of the outgoing fluxes will be entirely determined by the incoming fluxes.

We now formulate this question in mathematical terms: the incoming fluxes are described by two functions  $\mathcal{F}_0(p)$  (defined for  $p \geq 0$ ) and  $\mathcal{F}_1(p)$  (defined for  $p \leq 0$ ). These are the incoming distributions on the left end (index 0) and the right end (index 1) of the system.

In terms of these 2 functions, the problem of existence of a stationary state can be formulated as (recall that the rescaled system has length one):

*Is there a solution  $(F, g)$  of Eqs. (4.5) with the boundary conditions*

$$F(p, 0) = \mathcal{F}_0(p), \quad \forall p \geq 0 \quad \text{and} \quad F(p, 1) = \mathcal{F}_1(p), \quad \forall p \leq 0. \quad (5.1)$$

Assume for a moment that, instead of the boundary conditions (5.1) we were given just  $F(p, 0)$ , but now for all  $p \in \mathbb{R}$ , not only for  $p > 0$ . Assume furthermore, that  $g(p, x)$  is determined by (4.5b). In that case, the relation (4.5) can be written as a dynamical system in the variable  $x$ :

$$\partial_x F(\cdot, x) = \mathcal{X}(F(\cdot, x)). \quad (5.2)$$

Thus, if  $F(\cdot, 0)$  is given, then, in principle,  $F(\cdot, 1)$  is determined (uniquely) by Eq. (5.2), provided such a solution exists. We denote this map by  $\mathcal{B}_0$ :

$$\mathcal{B}_0 : F(\cdot, 0) \mapsto F(\cdot, 1).$$

What is of interest for our problem is the restriction of the image of  $\mathcal{B}$  to functions of negative  $p$  only, since that corresponds to the incoming particles from the right side, and so we define

$$(\mathcal{B}(F(\cdot, 0))) (p) \equiv \theta_-(p) \cdot (\mathcal{B}_0(F(\cdot, 0))) (p) = \theta_-(p) \cdot F(p, 1).$$

Using this map  $\mathcal{B}$ , we will show that when  $F(\cdot, 0)$  varies in a small neighborhood the map  $\mathcal{B}$  is invertible on its image. By taking inverses the problem of heat conduction for our model will be solved for small temperature and flux difference.

Of course, this needs a careful study of the function space on which  $\mathcal{B}$  is supposed to act. This will be done below.

To formulate the problem more precisely, we change notation, and let

$$\begin{aligned} F_0^+(p) &= \theta_+(p)F(p, 0), \\ F_0^-(p) &= \theta_-(p)F(p, 0), \\ F_1^+(p) &= \theta_+(p)F(p, 1), \\ F_1^-(p) &= \theta_-(p)F(p, 1). \end{aligned}$$

We assume now that  $F_0^+$  is *fixed* once and for all and omit it from the notation. Then, we see that  $\mathcal{B}$  can be interpreted as a map which maps the function  $F_0^-$  to  $F_1^-$ , and we call this map  $\Phi$ .

We will show below that for  $F_0^-$  in a small neighborhood  $\mathcal{D}$  the map  $\Phi$  is 1-1 onto its image  $\Phi(\mathcal{D})$  and can therefore be inverted. For any  $\hat{F}$  in  $\Phi(\mathcal{D})$ , we can take  $F_0^- = \Phi^{-1}(\hat{F})$ , and we will have solved the problem of existence of heat flux.

### 6. The $g$ Equation

We start here with the study of existence of  $g$  for given  $F$ . Since (4.5b) does not couple different  $x$ , we can fix  $x$ . Equation (4.5b) is then equivalent to

$$g = \mathcal{A}_F(g),$$

where the operator  $\mathcal{A}_F$  acting on the function  $h$  is defined by

$$\mathcal{A}_F(h)(\tilde{P}, x) = \frac{\int dp F(\tilde{p}(p, \tilde{P}), x) h(P(p, \tilde{P}), x)}{\int dp F(p, x)},$$

(provided the denominator does not vanish). Note that for fixed  $\tilde{P}$  and  $p$ , we can solve the collision system (4.2) to find the corresponding  $P$  and  $\tilde{p}$ , namely

$$P = \frac{1}{\varrho} \tilde{P} - \frac{1+\varrho}{\varrho} p \quad \text{and} \quad \tilde{p} = \frac{1-\varrho}{\varrho} \tilde{P} - \frac{1}{\varrho} p.$$

The action of the operator  $\mathcal{A}_F$  can then be rewritten as

$$\mathcal{A}_F(h)(\tilde{P}, x) = \frac{\varrho^{-1} \int dp F(p, x) h\left(\tilde{P}/\varrho - (1+\varrho)p/\varrho, x\right)}{\int ds F(s, x)}. \tag{6.1}$$

In order to study this operator notice that it does not depend explicitly on  $x$ . It is convenient to study instead a family of operators indexed by functions  $\varphi$  of the momentum only. We define (assuming the integral of  $\varphi$  does not vanish)

$$(\mathcal{L}_\varphi \psi)(\tilde{P}) = \frac{\int dp \varphi(p) \psi\left(\tilde{P}/\varrho - (1+\varrho)p/\varrho\right)}{\varrho \int dp \varphi(p)}.$$

A final change of variables will be useful when we study  $\mathcal{L}_\varphi$ :

$$(\mathcal{L}_\varphi \psi)(p) = \frac{\int dq \varphi\left(\frac{p-\varrho q}{1+\varrho}\right) \psi(q)}{(1+\varrho) \int dq \varphi(q)}. \tag{6.2}$$

### 7. The Mathematical Setup and the Main Result

Having formulated the problem of existence of the stationary solution in general, we now fix the mathematical framework in which we can prove this existence. This framework, while quite general, depends nevertheless on a certain number of technical assumptions which we formulate now.

We fix once and for all the ratio  $\mu = m/M$  of the masses, and assume, for definiteness, that  $\mu \in (0, 1)$ . It seems that this condition is not really necessary, and probably

the condition  $m \neq M$  (and the masses non-zero) should work as well, but we have not pursued this.

We next describe a condition on the incoming distribution, called  $F$  in the earlier sections. The basic idea, inspired from the equilibrium calculations, is that  $F(p, x)$  should be close to

$$F_{\text{reference}}(p) = \exp(-\beta p^2 / (2m)) \equiv \exp(-\alpha p^2),$$

while the derived quantity  $g(p, x)$  should be close to

$$g_{\text{reference}}(p) = \exp(-\beta p^2 / (2M)) \equiv \exp(-\mu \alpha p^2).$$

Upon rescaling  $p$ , we may assume henceforth that  $\alpha = 1$ .

The operators of the earlier sections will now be described in spaces with weights

$$\mathcal{W}_\nu(p) = \exp(-\nu p^2),$$

where we will choose  $\nu = 1$  for the  $F$  and  $\nu = \mu$  for the  $g$ .

We recall that the operator  $\mathcal{L}_F$  in “flat” space is

$$(\mathcal{L}_F g)(p) = \frac{\int dq F\left(\frac{p-\varrho q}{1+\varrho}\right) g(q)}{(1+\varrho) \int dq F(q)}.$$

We then define the integral kernel in the space with weights  $\exp(-p^2)$  for  $F$  and  $\exp(-\mu p^2)$  for  $g$ , and write

$$F(p) = e^{-p^2} v(p), \quad g(p) = e^{-\mu p^2} u(p).$$

Here,  $\mu = m/M = (1-\varrho)/(1+\varrho)$ , as before. Expressed with  $u$  and  $v$  the operator  $\mathcal{L}_F$  takes the form

$$(K_v u)(p) = \frac{1}{(1+\varrho) \int dq e^{-q^2} v(q)} \cdot (L_v u)(p), \tag{7.1}$$

where

$$(L_v u)(p) = \int dq v\left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q) u(q),$$

and

$$K(p, q) = \mathcal{W}_1\left(\frac{p-\varrho q}{1+\varrho}\right) \cdot \mathcal{W}_\mu(q) / \mathcal{W}_\mu(p).$$

A simple calculation shows that

$$K(p, q) = e^{-(\varrho p - q)^2 / (1+\varrho)^2}. \tag{7.2}$$

Our task will be to understand under which conditions the linear operator  $\mathcal{L}_F$  has an eigenvalue 1. This will be done by showing that  $K_v$  is quasi-compact. It is here that we were not able to give reasonable bounds on  $K(p, q)$  in the case of different exponentials for  $p > 0$  and  $p < 0$ , which represents different temperatures for ingoing and outgoing particles.

7.1. *Function spaces.* We now define spaces which are adapted to the simultaneous requirement of functions being close to a Gaussian near  $|p| = \infty$  and  $u$  and  $v$  having limits, and  $K_v$  being quasi-compact. We define a space  $\mathcal{G}_1$  of functions  $u$  with norm

$$\|u\|_{\mathcal{G}_1} = \int e^{-\mu p^2} |u(p)| dp.$$

Similarly,  $\mathcal{F}_1$  is the space of functions  $v$  with norm

$$\|v\|_{\mathcal{F}_1} = \int e^{-p^2} |v(p)| dp.$$

Thus, the only difference is the absence of the factor  $\mu = (1 - \rho)/(1 + \rho)$  in the exponent.

We also define a smaller space  $\mathcal{G}_2$ , contained in  $\mathcal{G}_1$ , with the norm

$$\|u\|_{\mathcal{G}_2} = \int |du(p)| + \int e^{-\mu p^2} |u(p)| dp,$$

and the analogous space  $\mathcal{F}_2$  contained in  $\mathcal{F}_1$  with the norm

$$\|v\|_{\mathcal{F}_2} = \int |dv(p)| + \int e^{-p^2} |v(p)| dp.$$

*Remark.* To simplify notation we write  $\int |du(p)|$  instead of the variation norm. However, the “integration by parts” formula would hold with the “correct” definition of variation as well.

**Lemma 7.1.** *One has the inclusion  $\mathcal{G}_2 \subset L^\infty$ , and more precisely*

$$\|u\|_{L^\infty} \leq \int |du| + e^\mu \int e^{-\mu p^2} |u(p)| dp \leq e^\mu \|u\|_{\mathcal{G}_2}.$$

Furthermore, if  $u \in \mathcal{G}_2$ , then  $\lim_{p \rightarrow \pm\infty} u(p)$  exists. The maps  $u \mapsto \lim_{p \rightarrow \pm\infty} u(p)$  and  $u \mapsto \int dp \exp(-\mu p^2) \cdot u(p)$  are continuous functions from  $\mathcal{G}_2$  to  $\mathbb{R}$ . The unit ball of  $\mathcal{G}_2$  is compact in  $\mathcal{G}_1$ .

Analogous statements hold for the spaces  $\mathcal{F}_2$  (defined without the factor  $\mu$ ).

*Proof.* The first statement is easy, but it will be convenient to have the explicit estimates. We have

$$u(y) - u(x) = \int_{[x,y]} du,$$

and therefore

$$|u(x)| \leq \int |du| + \int_{-1/2}^{1/2} |u(y)| dy \leq \int |du| + e^\mu \int e^{-\mu p^2} |u(p)| dp.$$

The second statement follows at once since the functions in  $\mathcal{G}_2$  are of bounded variation.

For the last assertions, it follows from the inclusion in  $L^\infty$  that the unit ball of  $\mathcal{G}_2$  is equi-integrable at infinity in  $L^1(e^{-\mu p^2} dp)$ . Moreover, a set of uniformly bounded functions of uniformly bounded variation is compact in any  $L^1(K, dp)$  for any compact subset  $K$  of  $\mathbb{R}$  (see [2], Helly’s selection principle).  $\square$

7.2. *A cone in  $\mathcal{F}_2$ .* We will work in the space  $\mathcal{F}_2$  but we will need a cone (of positive functions, with adequate decay) in this space, in order to prove quasi-compactness of  $K_v$ .

We define a cone  $\mathcal{C}_{\mathcal{F}}$  in  $\mathcal{F}_2$  by the condition

$$\mathcal{C}_{\mathcal{F}} = \left\{ v \in \mathcal{F}_2, v \geq 0 \text{ and } Z \cdot \lim_{p \rightarrow \pm\infty} v(p) < 1 \right\}, \tag{7.3}$$

where

$$Z = Z(v) = \frac{\sqrt{\pi}}{\int dp e^{-p^2} v(p)}. \tag{7.4}$$

**Lemma 7.2.** *The cone  $\mathcal{C}_{\mathcal{F}}$  has non empty interior (in  $\mathcal{F}_2$ ) and is convex.*

*Proof.* By Lemma 7.1 the maps  $v \mapsto \lim_{p \rightarrow \pm\infty} v(p)$  and  $v \mapsto \int \exp(-p^2)v(p)dp$  are continuous in  $\mathcal{F}_2$  and hence the assertion follows.  $\square$

*Remark.* Note that a function in the interior of the cone is necessarily bounded away from zero, since at infinity it must have a non-zero limit and in any compact set, if it is never zero, it is bounded away from zero.

*Remark.* Note that the function  $v \equiv 1$  (the Gaussian) is *not* in the cone  $\mathcal{C}_{\mathcal{F}}$ . In fact, we require that  $\lim_{p \rightarrow \pm\infty} F(p)e^{p^2} \cdot \int e^{-p'^2} dp' / \int F(p')dp' < 1$ .

7.3. *The main result.* On the set  $\mathcal{C}_{\mathcal{F}}$ , we consider now the spatial evolution equations (4.5) in the variables  $v$  and  $u_v$  (which is the solution of  $K_v u = u$  with  $K_v$  defined in (7.1)):

$$\begin{aligned} & \partial_x v(p, x) \\ &= \text{sign}(p) \int dP \left( (\mathcal{W}_1 \cdot v)(\tilde{p}, x) (\mathcal{W}_\mu \cdot u_{v(\cdot, x)})(\tilde{P}, x) \right. \\ & \quad \left. \times (\mathcal{W}_1 \cdot v)(p, x) (\mathcal{W}_\mu \cdot u_{v(\cdot, x)})(P, x) \right) \\ &= \text{sign}(p) \int dP (\mathcal{W}_1 \cdot v)(-qp + (1 - q)P, x) (\mathcal{W}_\mu \cdot u_{v(\cdot, x)})((1 + q)p + qP, x) \\ & \quad - \text{sign}(p) v(p, x) \int dP e^{-\mu P^2} u_{v(\cdot, x)}(P, x), \end{aligned} \tag{7.5}$$

with initial condition  $v(\cdot, x = 0) \in \mathcal{C}_{\mathcal{F}}$ . We will give a more explicit variant in (8.9).

Any solution of this equation is a function of  $p$  and  $x$ , and it is easy to verify that it satisfies Eq. (4.5a). Together with the definition of  $u_v$  we have a complete solution of the nonlinear system (4.5). Here we assume of course that the r.h.s. of the above equation is well defined as a function, so that we can multiply by  $\text{sign}(p)$ .

**Theorem 7.3.** *For any  $v_0 \in \mathcal{C}_{\mathcal{F}}$ , there are a number  $x_{v_0} > 0$  and a neighborhood  $\mathcal{V}_{v_0}$  of  $v_0$  in  $\mathcal{C}_{\mathcal{F}}$  such that the solution of (7.5) exists for any initial condition  $v_0 = v(p, 0) \in \mathcal{V}_{v_0}$  and for any  $x$  in the interval  $[0, x_{v_0}]$ . The function  $v_0 \mapsto x_{v_0}$  is continuous from  $\mathcal{C}_{\mathcal{F}}$  to  $\mathbb{R}^+$  compactified at infinity. We denote by  $\Phi_x$  the semi-flow integrating (7.5). For any  $x \in [0, x_{v_0}]$ , the map  $\Phi_x : v \mapsto \Phi_x(v)$  is a local diffeomorphism, i.e., a diffeomorphism on  $\mathcal{V}_{v_0}$ .*

Note that this implies in particular that the probability densities for  $g(\cdot, x)$  and  $F(\cdot, x)$  remain positive for all  $x \in [0, x_{v_0}]$ , which is of course crucial from the physics point of view.

We will prove this in Sect. 8 (and in the Appendix).

### 8. Bound on the Operator $K_v$ and Proof of Theorem 7.3

These bounds are the crux of the matter. They actually show, that, under the conditions on  $\mathcal{F}_2$  and the set  $\mathcal{C}_{\mathcal{F}}$ , the operator  $K_v$  is quasi-compact. In terms of the physical problem, this means that the scatterer is not heating up if the incoming fluxes are in  $\mathcal{F}_2$ .

The object of study is, for  $v \in \mathcal{C}_{\mathcal{F}}$ , the operator

$$(K_v u)(p) = \frac{1}{(1 + \varrho) \int e^{-q^2} v(q) dq} \int v\left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q) u(q) dq,$$

and we are asking for a solution  $u$  of the equation  $K_v(u) = u$ .

**Lemma 8.1.** *If  $v \geq 0$ ,  $K_v$  is a positive (nonnegative) operator, and*

$$\int dp e^{-\mu p^2} (K_v u)(p) = \int dp e^{-\mu p^2} u(p)$$

and

$$\|K_v\|_{\mathcal{G}_1} = 1.$$

*Proof.* Easy, compute and take absolute values. Alternately, consider that the probability is conserved in the original space.  $\square$

Our main technical result is

**Proposition 8.2.** *For  $v \in \mathcal{C}_{\mathcal{F}}$ , there exist a  $\zeta$ ,  $0 \leq \zeta < 1$  and an  $R > 0$  (both depend on  $v$  continuously) such that for any  $u \in \mathcal{G}_2$  one has the bound*

$$\int |dK_v(u)| \leq \zeta \int |du| + R \|u\|_{\mathcal{G}_1}.$$

*Proof.* Since  $v$  will be fixed throughout the study of  $K_v$ , it will be useful to introduce the abbreviation  $Q = Q_v$  for the normalizing factor

$$Q = \frac{1}{(1 + \varrho) \int e^{-q^2} v(q) dq}. \tag{8.1}$$

We will use a family of smooth cut-off functions  $\chi_L$  ( $L > 1$ ) which are equal to 1 on  $[-L + \frac{1}{2}, L - \frac{1}{2}]$  and which vanish on  $|q| > L + \frac{1}{2}$ . Let  $\Theta$  be a  $C^\infty$  function satisfying  $0 \leq \Theta \leq 1$ , with  $\Theta(q) = 0$  for  $q \leq -\frac{1}{2}$  and  $\Theta(q) = 1$  for  $q \geq \frac{1}{2}$ . We define  $\chi_L$  by

$$\chi_L(q) = \begin{cases} \Theta(q + L) & \text{if } q \leq -L + \frac{1}{2}, \\ 1 & \text{if } -L + \frac{1}{2} \leq q \leq L - \frac{1}{2}, \\ \Theta(L - q) & \text{if } q \geq L - \frac{1}{2}. \end{cases}$$



The functions  $\chi_L$  are  $C^\infty$ , satisfy  $0 \leq \chi_L \leq 1$  and  $\|\chi'_L\|_{L^\infty}$  is independent of  $L$ . Let  $L_1$  and  $L_2$  be two positive numbers to be chosen large enough later on (depending on  $v$ ). We will use the partition of unity

$$1 = \chi_L + \chi_L^\perp.$$

Using this decomposition of unity with  $L = L_1$  and  $L = L_2$ , we write  $K_v = K_v^{(1)} + K_v^{(2)} + K_v^{(3)}$  with

$$\begin{aligned} (K_v^{(1)}u)(p) &= Q \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) K(p, q)u(q) \cdot \chi_{L_1}(q), \\ (K_v^{(2)}u)(p) &= Q \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q), \\ (K_v^{(3)}u)(p) &= Q \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q). \end{aligned}$$

We will now estimate the variation of the three operators separately. For the variation of the first term, we find

$$\begin{aligned} d(K_v^{(1)}u)(p) &= +Q \int dq dv \left( \frac{p-\varrho q}{1+\varrho} \right) \frac{1}{1+\varrho} \cdot K(p, q)u(q) \chi_{L_1}(q) \\ &\quad + Q dp \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) \partial_p K(p, q) \cdot u(q) \chi_{L_1}(q). \end{aligned}$$

Using the explicit form of  $K(p, q)$  (see Eq. (7.2)), and some easy bounds which we defer to the Appendix, we get the bound

$$\begin{aligned} \int |d(K_v^{(1)}u)(p)| &\leq \mathcal{O}(1)Q \left( \|v\|_{L^\infty} + \int |dv| \right) \int_{-L_1-\frac{1}{2}}^{L_1+\frac{1}{2}} |u(q)|dq \\ &\leq \mathcal{O}(1)Q \left( \|v\|_{L^\infty} + \int |dv| \right) e^{\mu(L_1+\frac{1}{2})^2} \int e^{-\mu q^2} |u(q)|dq \quad (8.2) \\ &\leq \text{const.} \|u\|_{\mathcal{G}_1} \cdot \|v\|_{\mathcal{F}_2}. \end{aligned}$$

The variation of  $K_v^{(2)}$  leads to three terms:

$$\begin{aligned} d(K_v^{(2)}u)(p) &= Q dp \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi'_{L_2}(p - \varrho q) \\ &\quad + Q \int dq dv \left( \frac{p-\varrho q}{1+\varrho} \right) \frac{1}{1+\varrho} \cdot K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q) \\ &\quad + Q dp \int dq v \left( \frac{p-\varrho q}{1+\varrho} \right) \partial_p K(p, q) \cdot u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q) \\ &:= dJ_{21} + dJ_{22} + dJ_{23}. \end{aligned}$$

In these terms, the variables  $p$  and  $q$  are in the domain

$$\mathcal{D} = \{(p, q) \in \mathbb{R}^2 : |p - \varrho q| < L_2 + \frac{1}{2} \text{ and } |q| > L_1 - \frac{1}{2}\},$$

and for  $L_1 = 3\varrho L_2/(1 - \varrho^2)$  and  $L_2$  sufficiently large we have from Lemma A.3:

$$K(p, q) < \exp\left(-C_1(\varrho p - q)^2 - C_2 L_2^2\right). \tag{8.3}$$

Therefore, we get for  $dJ_{21}$ :

$$\int |dJ_{21}| \leq \text{const.} e^{-C_2 L_2^2} \|u\|_\infty \|v\|_\infty \int_{(p,q) \in \mathcal{D}} dp dq e^{-C_1(\varrho p - q)^2}. \tag{8.4}$$

The integral exists and is uniformly bounded in  $L_2$  (since  $|\varrho p - q| \rightarrow \infty$  when  $|q| \rightarrow \infty$ ). The term  $dJ_{22}$  is handled in a similar way and leads to the bound

$$\int |dJ_{22}| \leq \text{const.} e^{-C_2 L_2^2} \|u\|_\infty \int |dv|. \tag{8.5}$$

The following identity is useful:

$$\partial_p K(p, q) = -\varrho \partial_q K(p, q). \tag{8.6}$$

For the term  $dJ_{23}$  we observe that from (8.6) one gets, upon integrating by parts, with the notation

$$\begin{aligned} dJ_{23} &= \varrho Q dp \int dq v \left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q) \cdot \partial_q \left(\chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q)\right) \\ &\quad + \varrho Q dp \int dv \left(\frac{p-\varrho q}{1+\varrho}\right) \frac{-\varrho}{1+\varrho} K(p, q) u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q) \\ &\quad + \varrho Q dp \int v \left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q) du(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q) \\ &:= dJ_{231} + dJ_{232} + dJ_{233}. \end{aligned}$$

All these terms are localized in the domain  $\mathcal{D}$ . In  $dJ_{231}$  there appears a derivative

$$\begin{aligned} X &= \partial_q \left(\chi_{L_1}^\perp(q) \chi_{L_2}(p - \varrho q)\right) \\ &= -\chi'_{L_1}(q) \chi_{L_2}(p - \varrho q) \\ &\quad - \varrho \chi_{L_1}^\perp(q) \chi'_{L_2}(p - \varrho q) \\ &:= X_1 + X_2. \end{aligned}$$

The terms involving  $X_1$  and  $X_2$  can be bounded as  $dJ_{21}$  and  $dJ_{22}$  by observing that  $\text{supp} \chi'_{L_1} \subset \{|q| < L_1 + \frac{1}{2}\}$ , and similarly for  $X_2$ .

The terms  $dJ_{232}$  and  $dJ_{233}$  are bounded similarly.

Together, these lead to a bound

$$\int |dK_v^{(2)}(u)| \leq \text{const.} e^{-C_2 L_2^2} \|v\|_{\mathcal{F}_2} \|u\|_{\mathcal{G}_2}. \tag{8.7}$$

*Remark.* Note that in this term, the norm  $\|u\|_{\mathcal{G}_2}$  appears with a *small* coefficient, while in (8.2) it was  $\|u\|_{\mathcal{G}_1}$  (with a large coefficient).

Finally, we estimate the total variation of  $K_v^{(3)}(u)$  and here, the nature of the set  $\mathcal{C}_{\mathcal{F}}$  will be important. We have

$$\begin{aligned} & d\left(K_v^{(3)}u\right)(p) \\ &= Q \int dq dv \left(\frac{p-\varrho q}{1+\varrho}\right) \frac{1}{1+\varrho} K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q) \\ &\quad - Q dp \int dq v \left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi'_{L_2}(p - \varrho q) \\ &\quad + Q dp \int dq v \left(\frac{p-\varrho q}{1+\varrho}\right) \partial_p K(p, q) \cdot u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q) \\ &:= dJ_{31} + dJ_{32} + dJ_{33}. \end{aligned}$$

The critical term is  $dJ_{33}$ , but we first deal with the two others which are treated similar to earlier cases.

For the first term we have by Lemma A.4 which tells us that  $K$  is exponentially bounded on  $\mathcal{D}'$ :

$$\int |dJ_{31}| \leq \text{const.} \|u\|_{L^\infty} \int_{|s| > (L_2 - \frac{1}{2})/(1+\varrho)} |dv(s)|,$$

where  $\mathcal{D}'$  is the domain

$$\mathcal{D}' = \{(p, q) : |q| > L_1 \quad \text{and} \quad |p - \varrho q| > L_2\}.$$

For the second term, we have, again by Lemma A.4 below,

$$\int |dJ_{32}| \leq \text{const.} \|u\|_{L^\infty} \|v\|_{L^\infty}.$$

The last term is more delicate, and uses the property  $Z \cdot \lim_{p \rightarrow \pm\infty} v(p) < 1$  of the definition of the cone  $\mathcal{C}_{\mathcal{F}}$ , Eq. (7.3). integrate by parts as before using (8.6) and get

$$\begin{aligned} dJ_{33} &= \varrho Q dp \int dv \left(\frac{p-\varrho q}{1+\varrho}\right) \frac{-\varrho}{1+\varrho} K(p, q)u(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q) \\ &\quad \times \varrho Q dp \int dq v \left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q)u(q) \cdot \partial_q \left(\chi_{L_1}(q) \chi_{L_2}^\perp(p - \varrho q)\right) \\ &\quad \times \varrho Q dp \int v \left(\frac{p-\varrho q}{1+\varrho}\right) K(p, q)du(q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q) \\ &:= dJ_{331} + dJ_{332} + dJ_{333}. \end{aligned}$$

The term  $dJ_{331}$  is bounded like  $dJ_{31}$ .

In a similar way  $dJ_{332}$  and  $dJ_{333}$  are bounded by the same methods.

The term  $dJ_{333}$  makes use of the limit condition in  $\mathcal{C}_{\mathcal{F}}$ . Consider the integral of  $|dJ_{333}|$ . This leads to a bound and setting  $L'_2 = (L_2 - \frac{1}{2})/(1 + \varrho)$ :

$$\begin{aligned} \int |dJ_{333}(p)| &\leq \varrho Q \sup_{|s|>L'_2} |v(s)| \cdot \int |du(q)| \\ &\quad \cdot \int dp K(p, q) \cdot \chi_{L_1}^\perp(q) \chi_{L_2}^\perp(p - \varrho q) \\ &\leq \varrho Q \left( \sup_{|q|>L_1 - \frac{1}{2}} \int dp K(p, q) \right) \cdot \sup_{|s|>L'_2} |v(s)| \int |du| \quad (8.8) \\ &= \frac{\sqrt{\pi}}{\int dq e^{-q^2} v(q)} \cdot \sup_{|s|>L'_2} |v(s)| \int |du| \\ &= Z \cdot \sup_{|s|>L'_2} |v(s)| \int |du|, \end{aligned}$$

where  $Z$  was defined in Eq. (7.4). Collecting all the estimates, we get

$$\int |dK_v(u)| \leq C \int e^{-\mu q^2} |u(q)| dq + \zeta(L_2) \int |du|,$$

where

$$\zeta(L_2) = \mathcal{O}(1)e^{-C_2 L_2^2} \|v\|_{\mathcal{F}_2} + \mathcal{O}(1) \int_{|s|>L'_2} |dv(s)| + Z \cdot \sup_{|s|>L'_2} |v(s)|.$$

Since  $v$  belongs to  $\mathcal{C}_{\mathcal{F}}$ , it follows that

$$\lim_{L_2 \rightarrow \infty} \zeta(L_2) < 1,$$

and the lemma follows by taking  $L_2$  large enough.  $\square$

**Proposition 8.3.** *For any  $v \in \mathcal{C}_{\mathcal{F}}$ , the equation  $K_v(u) = u$  has a solution in  $\mathcal{G}_2$ . This solution can be chosen positive, it is then unique if we impose  $\|u\|_{\mathcal{G}_1} = 1$ . We call it  $u_v$ . The map  $v \mapsto u_v$  is differentiable.*

*Proof.* We apply the theorem of Ionescu-Tulcea and Marinescu [5] to prove the existence of  $u$ . Since for  $v > 0$ , the operator  $K_v$  is positivity improving, it follows by a well known argument, see e.g., [7] that the peripheral spectrum consists only of the simple eigenvalue one and the eigenvector can be chosen positive. If normalized, it is then unique. Since the operator  $K_v$  depends linearly and continuously on  $v$  (in  $\mathcal{F}_2$ ), the last result follows by analytic perturbation theory (see [6]).  $\square$

We next consider Eq. (7.5) for  $v$ :

$$\begin{aligned} \partial_x v(p) &= \text{sign}(p) \left( \frac{1}{\varrho} \int e^{(1-\varrho^2)(p-q)/(1+\varrho))^2/\varrho^2} v\left(\frac{(1-\varrho)q-p}{\varrho}\right) u_v(q) dq \right. \\ &\quad \left. - v(p) \int e^{-\mu q^2} u_v(q) dq \right). \quad (8.9) \end{aligned}$$

**Proposition 8.4.** *The r.h.s. of the equation for  $v$  is a  $C^1$  vector field on  $\mathcal{F}_2$ .*

*Proof.* This follows easily from the fact that the map  $v \mapsto u_v$  is  $C^1$ .  $\square$

**Theorem 8.5.** *Let  $v_0 \in \mathcal{C}_{\mathcal{F}}$ , and assume that  $v_0$  is bounded below away from zero and has nonzero limits at  $\pm\infty$ . Then there is a number  $s = s(v) > 0$  such that the solution Eq. (8.9) with initial condition  $v_0$  exists in  $\mathcal{F}_2$  and is nonnegative (moreover, it belongs to  $\mathcal{C}_{\mathcal{F}}$ ).*

*Proof.* Follows at once from the previous proposition and the fact that  $v_0$  is in the interior of  $\mathcal{C}_{\mathcal{F}}$ .  $\square$

The proof of Theorem 7.3 is now completed by observing that the map  $\Phi : v_0 \mapsto \Phi(v_0)$  is indeed a local diffeomorphism, since it is given as the solution of an evolution equation.

### 9. Remarks and Discussion

9.1. *The behavior of the solution at  $p = \infty$ .* Consider the limit  $p \rightarrow \infty$  in the expression for  $K_v$ . We need  $\varrho p - q = \mathcal{O}(1)$  otherwise the Gaussian gives a negligible contribution. In other words,  $q \sim \varrho p$ , and we are going to assume from now on that  $\varrho > 0$  (the other case can be treated analogously). This implies  $p - \varrho q \sim (1 - \varrho^2)p$  which also tends to infinity (the same infinity). Therefore,

$$K_v u(\pm\infty) = \frac{\sqrt{\pi} v(\pm\infty) u(\pm\infty)}{\int e^{-p^2} v(p) dp}.$$

In particular, if  $K_v u = u$  and since we assumed

$$\frac{\sqrt{\pi} v(\pm\infty)}{\int e^{-p^2} v(p) dp} \neq 1$$

we get  $u(\pm\infty) = 0$ .

For the  $v$  equation, we have for large  $p$ ,  $q \sim p(1 + \varrho)$  and  $(1 - \varrho)q - p \sim -\varrho^2 p$ . Therefore (inverting limit and derivative) we get

$$\partial_x v(\pm\infty) = \text{sign}(\pm\infty) \left[ \sqrt{\pi} \sqrt{\frac{1+\varrho}{1-\varrho}} u(\pm\infty) v(\mp\infty) - v(\pm\infty) \int e^{-\mu q^2} u(q) dq \right].$$

Note that the first term vanishes since  $u(\pm\infty) = 0$ . Since the integral  $C(x) = \int e^{-\mu q^2} u(q, x)$  is positive, we conclude that formally,

$$\partial_x v(\pm\infty) = \mp v(\pm\infty) C(x).$$

9.2. Essential spectrum.

**Conjecture.** *The essential spectrum of  $K_v$  is the interval  $[0, \sigma(v)]$  with*

$$\sigma(v) = \max \frac{\sqrt{\pi} v(\pm\infty)}{\int e^{-p^2} v(p) dp}.$$

If  $\sigma(v) < 1$  we are looking for an eigenvalue 1 outside the essential spectrum, which is the case we have treated. If  $\sigma(v) > 1$  we would be looking for an eigenvalue 1 inside the essential spectrum which would be a much more difficult task, since it may well not exist.

*Idea of proof.* Similar to the above estimates, the operator  $K_v$  should be written as something small plus something compact plus something whose essential spectrum can be computed. This last part is likely to be the limit operator at infinity.

9.3. Dependence on  $N$ . It should be noted that the equation for  $\partial_x F$  has, in fact a scaling of the form

$$N^{-1} \partial_x F = \mathcal{O}(1) + \mathcal{O}(N^{-1}).$$

This means that in the main theorem (Theorem 7.3), the limit  $x_{v_0}$  of  $x$  for which we have a result is quite probably bounded by a quantity of the form  $1/(N \cdot \Delta(v_0))$ , where  $\Delta(v_0)$  measures the deviation of the initial condition  $v_0$  from a Gaussian. Thus, either  $x_{v_0}$  is very small when  $N$  is large, or one has to take  $v_0$  very close to a Gaussian.

Another way to look at this scaling is to introduce a scattering probability  $\gamma = b/N$  where  $b > 0$  is a constant independent of  $N$ . In other words, a particle entering the array of cells from the left has for large  $N$  a probability  $e^{-b}$  to traverse all the  $N$  cells (and leave on the right) without having experienced any scattering. This is analogous to a rarefied gas. It is easy to verify that Eq. (2.4) is modified by a factor  $b/N$  multiplying the right hand side, and hence Eq. (2.5) is unchanged. The stationary equations (4.5) become

$$\begin{aligned} \frac{|\tilde{p}|}{m} f_L^-(t, \tilde{p}) &= \theta(-\tilde{p}) \left(1 - \frac{b}{N}\right) \frac{|\tilde{p}|}{m} f(t, \tilde{p}) \\ &+ \frac{b}{N} \frac{\theta(-\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g \left(t - mL/|\tilde{p}|, \frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f \left(t - \frac{m}{|\tilde{p}|} L - \frac{m}{|p|} L, p\right), \end{aligned}$$

and

$$\begin{aligned} \frac{|\tilde{p}|}{m} f_R^+(t, \tilde{p}) &= \theta(+\tilde{p}) \left(1 - \frac{b}{N}\right) \frac{|\tilde{p}|}{m} f(t, \tilde{p}) \\ &+ \frac{b}{N} \frac{\theta(+\tilde{p})}{1 - \varrho} \int_{\mathbb{R}} dp g \left(t - mL/|\tilde{p}|, \frac{\tilde{p} + \varrho p}{1 - \varrho}\right) \frac{|p|}{m} f \left(t - \frac{m}{|\tilde{p}|} L - \frac{m}{|p|} L, p\right). \end{aligned}$$

Equation (4.5a) follows as explained in Sect. 4 after a rescaling of space by a factor  $b$ .

9.4. *Discussion.* The model presented in this paper has the nice property that one can control the existence of a solution out of equilibrium. In particular, this means that there is no heating up of the scatterers in the “chain”, when the system is out of equilibrium.

The reader should note, however, that the initial condition at the boundary does not allow for different temperatures in the strict sense, only for different distributions at the ends. For example, a function of the form

$$F(p, 0) = \begin{cases} \exp(-\alpha p^2), & \text{if } p > 0, \\ \exp(-\alpha' p^2), & \text{if } p < 0, \end{cases}$$

with  $\alpha \neq \alpha'$  is not covered by Theorem 7.3. The reason for this failure is that we could not find an adequate analog of Lemma A.4 for initial conditions of this type, and therefore the bounds on the kernel  $K(p, q)$  are not good enough.

*Acknowledgement.* We thank Ph. Jacquet for a careful reading of the manuscript. This work was partially supported by the Fonds National Suisse.

**A. Appendix: Bounds on  $K(p, q)$**

We study here the kernel  $K$  of (7.2), which equals

$$K(p, q) = e^{E(p, q)},$$

with

$$E(p, q) = \mu p^2 - \mu q^2 - \left(\frac{p - \varrho q}{1 + \varrho}\right)^2 = -(\varrho p - q)^2 / (1 + \varrho)^2. \tag{A.1}$$

**Lemma A.1.** *Assume  $|q| < L$ . There are constants  $C = C(L, \varrho)$  and  $D = D(L, \varrho) > 0$  such that for all  $p$ ,*

$$K(p, q) < C e^{-Dp^2}, \tag{A.2}$$

and

$$|\partial_p K(p, q)| < C e^{-Dp^2}. \tag{A.3}$$

*Proof.* Obvious.  $\square$

**Lemma A.2.** *Assume  $|p - \varrho q| < L$ . There are constants  $C = C(L, \varrho)$  and  $D = D(L, \varrho) > 0$  such that for all  $q$ ,*

$$K(p, q) < C e^{-Dq^2}, \tag{A.4}$$

and

$$|\partial_p K(p, q)| < C e^{-Dq^2}, \tag{A.5}$$

*Proof.* The proof is as in Lemma A.1, with the difference that now  $|p - \varrho q| < L$ .  $\square$

**Lemma A.3.** Consider the domain  $\mathcal{D}$  defined by

$$\mathcal{D} = \{(p, q) \in \mathbb{R}^2 : |p - \varrho q| < L_2 + \frac{1}{2} \text{ and } |q| > L_1 - \frac{1}{2}\}, \tag{A.6}$$

with

$$L_1 = \frac{3\varrho}{1 - \varrho^2} L_2. \tag{A.7}$$

For fixed  $\varrho \in (0, 1)$  and sufficiently large  $L_2$  there are positive constants  $C_1$  and  $C_2$  such that for  $(p, q) \in \mathcal{D}$  one has the bound

$$K(p, q) < \exp\left(-C_1(\varrho p - q)^2 - C_2 L_2^2\right).$$

*Proof.* From the definition of  $\mathcal{D}$  and  $(1 - \varrho^2)q = (\varrho p - q) - \varrho(p - \varrho q)$ , we find (for sufficiently large  $L_2$ ):

$$|\varrho p - q| \geq (1 - \varrho^2) |q| - \varrho |p - \varrho q| \geq (1 - \varrho^2) \left(L_1 - \frac{1}{2}\right) - \varrho \left(L_2 + \frac{1}{2}\right) > \varrho L_2. \tag{A.8}$$

Using the form

$$(\varrho p - q)^2 > \frac{1}{4}(\varrho p - q)^2 + \frac{1}{4}\varrho^2 L_2^2, \tag{A.9}$$

the assertion follows immediately.  $\square$

We next study the region

$$\mathcal{D}' = \{(p, q) : |q| > L_1 \text{ and } |p - \varrho q| > L_2\}. \tag{A.10}$$

In this region, we have the obvious bound

**Lemma A.4.** For  $(p, q) \in \mathcal{D}'$ , one has the bound

$$E(p, q) = -\left(\frac{\varrho p - q}{1 + \varrho}\right)^2.$$

### References

1. Collet, P., Eckmann, J.-P., Mejía-Monasterio, C.: ArXiv:0810.4464
2. Dunford, N., Schwartz, J.T.: *Linear operators. Part II: Spectral theory. Self adjoint operators in Hilbert space.* With the assistance of William G. Bade and Robert G. Bartle, New York-London: Interscience Publishers John Wiley & Sons, 1963
3. Eckmann, J.-P., Mejía-Monasterio, C., Zabey, E.: Memory effects in nonequilibrium transport for deterministic Hamiltonian systems. *J. Stat. Phys.* **123**, 1339–1360 (2006)
4. Eckmann, J.-P., Young, L.-S.: Nonequilibrium energy profiles for a class of 1-D models. *Commun. Math. Phys.* **262**, 237–267 (2006)
5. Ionescu Tulcea, C.T., Marinescu, G.: Théorie ergodique pour des classes d’opérations non complètement continues. *Ann. of Math. (2)* **52**, 140–147 (1950)
6. Kato, T.: *Perturbation Theory for Linear Operators.* Berlin: Springer-Verlag, 1984 (Second corrected printing of the second edition)
7. Schaefer, H.H., Wolff, M.P.: *Topological vector spaces.* second edition, Volume **3** of Graduate Texts in Mathematics, New York: Springer-Verlag, 1999