UNIQUENESS OF THE NON-EQUILIBRIUM STEADY STATE FOR A 1D BGK MODEL IN KINETIC THEORY

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Abstract. We continue our investigation of kinetic models of a one-dimensional gas in contact with homogeneous thermal reservoirs at different temperatures. Nonlinear collisional interactions between particles are modeled by a so-called BGK dynamics which conserves local energy and particle density. Weighting the nonlinear BGK term with a parameter $\alpha \in [0, 1]$, and the linear interaction with the reservoirs by $(1 - \alpha)$, we prove that for some $\alpha$ close enough to zero, the explicit spatially uniform non-equilibrium stable state (NESS) is unique, and there are no spatially non-uniform NESS with a spatial density $\rho$ belonging to $L^p$ for any $p > 1$. We also show that for all $\alpha \in [0, 1]$, the spatially uniform NESS is dynamically stable, with small perturbation converging to zero exponentially fast.

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1. Introduction

This paper is a contribution to the theory of non-equilibrium steady states (NESS), of open systems in the particular context of kinetic theory. The understanding of NESS, their properties, uniqueness or lack thereof and stability or lack thereof, represents a challenge in mathematical physics due to the fact that the dynamics are nonlinear, non-Markovian and the absence of an entropy principle. Our main result is a uniqueness and stability theorem for the NESS in a simple nonlinear model.

1.1. The model. We briefly describe the sort of underlying particle model that would lead to the type of kinetic equation that we study here. It consists of a gas of particles on the one-dimensional torus $\mathbb{T}$, that interact only through binary energy conserving collisions, however we also suppose that there are two types of scatterers distributed on the torus according to some Poisson distribution, as in a Lorentz
model, except that each scatterer has a temperature, $T_1$ or $T_2$ depending on its
type, and a certain radius of interaction, so that when a gas particle travels across
the interaction interval, a Poisson clock runs and if it goes off, the particle assumes
a new velocity chosen at random according to the Maxwellian distribution for the
temperature of the scatterer.

In an appropriate scaling limit, the net effect of the background scatterers is to
produce two uniform thermal reservoirs. Whatever the speed of a gas particle, its
rate of interaction with the reservoirs depends only on the density of the scatterers,
again in an appropriate limit in which their intervals of interaction are unlikely to
overlap. The kinetic equation that one would expect to arise from such a model
in such a limit would be of the type (1.1) below, except that one might expect a
Kac-Povzner type collision kernel, also known as a “soft-spheres” kernel [7, 13]. Our
work concerns the kinetic equation itself, and not its rigorous derivation from an
underlying particle system, although the brief description of such a system that we
have given hopefully illuminates the physical context of our model.

We are concerned with the existence and uniqueness of NESS for our system.
We make a further simplification, and model the gas particle collisions with a BGK
collision kernel [4, 14, 15]. This will render the existence of NESS trivial, but the
uniqueness is still a challenging problem, and we shall only prove part of what we
conjecture to be true.

These considerations bring us to the following one dimensional kinetic model:

$$
(1.1) \quad \partial_t f + v \partial_x f = \alpha M_f + (1 - \alpha) \rho_f \left( \frac{M_{T_1} + M_{T_2}}{2} \right) - f
$$

where $\alpha \in [0, 1]$, $f = f(t, x, v)$, $x \in \mathbb{T}$, $v \in \mathbb{R}$, and

$$
(1.2) \quad M_{T_i}(v) := \frac{e^{-\frac{|v|^2}{2T_i}}}{\sqrt{2\pi T_i}}, \quad M_f(t, x, v) := \rho_f(t, x) \frac{e^{-\frac{|v|^2}{2T_f(t, x)}}}{\sqrt{2\pi T_f(t, x)}},
$$

with

$$
(1.3) \quad \begin{cases}
\rho_f(t, x) := \int_{\mathbb{R}} f(t, x, v) \, dv, \\
P_f(t, x) := \int_{\mathbb{R}} v^2 f(t, x, v) \, dv = \rho_f(x) T_f(t, x),
\end{cases}
$$

being the spatial density and pressure corresponding to $f$. $T_f$ is then the temperature
corresponding to $f$, and $T_1$ and $T_2 \in (0, +\infty)$ being the two temperatures of the
reservoirs.

The linear terms on the right in (1.1) that are multiplied by $(1 - \alpha)$ model the
interaction of particles with two reservoirs, both acting everywhere in space. Each
time a particle interacts with one of the reservoirs, it velocity is replaced by a new
velocity selected at random from the corresponding Maxwellian distribution. We
have taken both of these Maxwellians to have zero mean velocity which is natural
for static reservoirs.
The nonlinear term on the right in (1.1) that is multiplied by $\alpha$ represents the effect of collisions between particles. The collision term $M_f$ is of the BGK type (see [4]), except that as usual in one dimensional kinetic models, it conserves only mass and energy, not momentum. Indeed, binary collisions that conserve both energy and momentum are trivial in one dimension: only an exchange of velocities is possible. For this reason, the mean velocity of $M_f$ is zero.

The term $-f$ on the right in (1.1) is the loss term corresponding to both interactions with the reservoirs and other particles: after such interactions, particles vacate their pre-interaction state.

Without loss of generality, we choose units in which the torus has unit volume and there is unit total mass:

$$T = [-1/2, 1/2] \text{ and } \int_{-1/2}^{1/2} \int_{\mathbb{R}} f(x, v) \, dx \, dv = 1.$$  

1.2. Previous results. In our previous papers [5, 6], we have studied related issues for related models. In [5] we proved the existence of spatially homogeneous non-equilibrium steady states and exponential convergence to them for related spatially homogeneous models, but with more realistic collision mechanisms, and also in higher dimensions. In [6] we studied the exponential rate of convergence to steady state for a non spatially homogeneous equation of the type (1.1) but with a modified collision mechanism that permitted the equation to be interpreted as the Kolmogorov forward equation for a non-stationary Markov process: we replaced the space-dependent local temperature $T_f(t, x)$ by the global temperature $T_f(t) := \int_{\mathbb{T} \times \mathbb{R}} v^2 f(x, v) \, dx \, dv$ of $f$ that depends only on time. We were then able to apply Doeblin’s method [16] to prove the exponential convergence. The use of Doeblin’s method to study linear dynamical models originates with [3, 11].

The rigorous study of NESS for nonlinear kinetic equations remains very challenging. One problem that has been studied by several authors is the Boltzmann equation in a slab with different temperatures on the two walls, with and without external forces. At this level of generality, one cannot always expect a unique NESS – there may be a symmetry breaking transition, such as the onset of Rayleigh-Bernard flow. Even without external forces, existence of NESS for the slab problem is a highly non-trivial, and existing results [2, 8] do not yield provide any information on uniqueness or non-uniqueness.

More recently, the Boltzmann equation in more general domains and with non-isothermal boundary conditions has been investigated [10] where it is proved that when the temperature on the boundary is sufficiently close to constant, then there is an NESS that is close to the uniform Maxwellian for the mean boundary temperature, and in a small neighborhood of this Maxwellian, there is no other NESS. However, it is not known that there are not other NESS further away, no matter how small the non-zero temperature difference may be.
1.3. Question studied and conjecture. Observe that the reservoirs will tend to damp out any mean velocity since \( M_{T_1} \) and \( M_{T_2} \) have zero mean velocity. Likewise, \( M_f \) has zero mean velocity at each \( x \), so the collision gain term too will tend to damp out any mean velocity.

Therefore, if \( f = f(v) \) is any spatially homogeneous steady state, \( \int_{\mathbb{R}} v f dv = 0 \). Moreover, the time and space homogeneity yield \( \partial_t f = 0 \) and \( v \partial_x f = 0 \). Finally, multiplying both sides of (1.1) by \( v^2 \) and integrating over \( x \) and \( v \) shows

\[
0 = (\alpha - 1)T_f + (1 - \alpha)\frac{T_1 + T_2}{2}, \quad T_f := \int_{\mathbb{R}} v^2 f(v) dv.
\]

Thus, the constant temperature in any spatially homogeneous steady state \( f \) must be \( T_\infty := (T_1 + T_2)/2 \) if \( \alpha \neq 1 \). Then for any spatially homogeneous steady state \( f \), \( M_f = M(T_1 + T_2)/2 \), and (1.1) reduces to

\[
\alpha M_{T_1 + T_2} + (1 - \alpha)\rho_f \left( \frac{M_{T_1} + M_{T_2}}{2} \right) - f = 0.
\]

Therefore the unique spatially homogeneous steady state is given by

\[
(1.4) \quad f_\infty := \alpha M_{T_\infty} + (1 - \alpha)\frac{M_{T_1} + M_{T_2}}{2}, \quad T_\infty = \frac{T_1 + T_2}{2}.
\]

Observe that \( f_\infty \) is not Maxwellian as soon as \( \alpha \neq 1 \).

If \( \alpha = 0 \), the term \( M_f \) is not present, the only spatially homogeneous steady state is \( f_\infty = \frac{1}{2}(M_{T_1} + M_{T_2}) \), and the equation (1.1) is linear. It can be interpreted as the forward equation of a Markov process and in [6] we used probabilistic methods to prove that this steady state is unique and is approached exponentially fast. Hence, for \( \alpha = 0 \), there are no steady states that are spatially inhomogeneous.

Next, consider the case \( \alpha = 1 \): there are no thermal reservoirs and energy is conserved. There is a one-parameter infinite family of steady states, namely \( M_T \) for all \( T > 0 \). Moreover, if \( f_0 \) is such that

\[
\int_{T \times \mathbb{R}} v^2 f_0(x,v) dx dv = T, \quad \int_{T \times \mathbb{R}} f_0(x,v) \ln f_0(x,v) dx dv < +\infty,
\]

and \( f(t,x,v) \) is the solution of (1.1) with initial datum \( f_0 \), then

\[
H(f(t,\cdot,\cdot),|M_T|) = \int_{T \times \mathbb{R}} f(t,x,v) \ln \frac{f(t,x,v)}{M_T(v)} dx dv
\]

decreases monotonically to zero, and is stationary only when \( f = M_T \). It follows that \( M_T \) is the unique steady state among solutions with second moment equal to \( T \) and finite entropy, and thus every steady state for \( \alpha = 1 \) with finite second moment and entropy is spatially homogeneous (and equal to \( M_T \)).

The question that motivates this paper is the study of the NESS in the intermediate region \( \alpha \in (0,1) \). We conjecture the following:

Conjecture (Uniqueness of the NESS for the BGK model with reservoirs). For all \( \alpha \in [0,1] \), the non-equilibrium steady state of (1.1) is unique, regardless of the temperature difference, spatially homogeneous and stable under perturbations. We also expect this conjecture to hold in higher dimensions \( x \in \mathbb{T}^d \), \( v \in \mathbb{R}^d \).
1.4. **Main results.** We give a partial answer to this conjecture, showing that it is satisfied when $\alpha$ is small enough. We first prove the uniqueness:

**Theorem 1.** For all $T_1, T_2$, there is an explicitly computable $\alpha_0 > 0$ such that for all $\alpha \in [0, \alpha_0)$, every steady state solution $f_{\infty}$ of \([1.1]\) that belongs to $L^1(\mathbb{T} \times \mathbb{R})$, has finite second moment and is such that $\rho \in L^p(\mathbb{T})$ for some $p > 1$, is constant in $x$.

**Remark 1.1.** Our method would also apply in higher dimensions provided there was non-trivial spatial dependence in only one direction on the torus, say the $x_1$ coordinate. The decomposition between odd and even parts in the next section should be then modified by splitting along $v_1$ only. The rest of the analysis would be similar.

We also prove the stability under perturbation for all $\alpha \in [0, 1]$. For this, we introduce the (real) Hilbert Space $\mathcal{H}_\alpha^1$ with inner product

$$\langle f, g \rangle_{\mathcal{H}_\alpha^1} = \int_{\mathbb{T} \times \mathbb{R}} (f(x, v)(1 - \partial^2_x)g(x, v)) \frac{1}{f_{\alpha, \infty}} \, dx \, dv .$$

**Theorem 2.** For all $\alpha \in [0, 1]$, the spatially homogeneous steady state described above is asymptotically stable under perturbation in $\mathcal{H}_\alpha^1$. Small perturbations decay exponentially fast in this space.

Theorem 2 shows that if for some $\alpha > \alpha_0$ there do exist non-uniform steady states, they do not arise as a branch bifurcating off the family of spatially homogeneous steady state solutions.

1.5. **Plan of the paper.** In Section 2, we establish some useful relations on the moments of any given NESS and introduce a decomposition between odd and even parts. In Section 3, we prove lower and upper pointwise bounds on local density and temperature of any given NESS. In Section 4, we explain the contraction mapping argument; it is in this section that we use that $\alpha$ is close to zero. Finally in Section 5, we prove a local stability result of the spatially homogeneous steady states for all $\alpha \in [0, 1]$.

2. **Preliminaries: properties and decompositions of NESS**

2.1. **Zero momentum and constant pressure.** A partial result supporting our conjecture is that the pressure is independent of $x$, as well as the momentum $m_f(x) := \rho_f(x)u_f(x)$, the latter being zero:

**Lemma 3.** Let $f(x, v)$ be a probability density on $\mathbb{T} \times \mathbb{R}$ such that $v^2f(x, v)$ is integrable, and suppose that $f(x, v)$ solves in a weak sense the equation

$$v\partial_x f(x, v) = F(x, v) - f(x, v)$$

where $F(x, v)$ is a measurable function such that $(1 + |v|)F(x, v)$ is integrable and

$$\forall x \in \mathbb{T}, \quad \int_{\mathbb{R}} F(x, v) \, dv = \rho_f(x) \quad \text{and} \quad \int_{\mathbb{R}} vF(x, v) \, dv = 0.$$
Then the pressure is constant and the momentum is zero:

\[
\begin{aligned}
P_f(x) &= \int_{\mathbb{R}} v^2 f(x, v) \, dv = P_\infty \in \mathbb{R} \\
m_f(x) &= \int_{\mathbb{R}} v f(x, v) \, dv = \rho_f(x) u_f(x) = 0.
\end{aligned}
\]

Remark 2.1. Evidently, Lemma 3 applies to any finite energy NESS of our equation.

Proof. Integrating both sides of (2.1) in \( v \) yields

\[
\frac{d}{dx} m_f(x) = \rho_f(x) - \rho_f(x) = 0.
\]

This proves that \( m_f(x) \) is a constant \( m_\infty \in \mathbb{R} \). Now multiplying both sides of (2.1) by \( v \) and integrating in \( v \) yields

\[
(2.3) \quad \frac{d}{dx} P_f(x) = -m_f(x) = -m_\infty.
\]

Integrating both sides of (2.3) in \( x \) shows that \( m_\infty = 0 \), and \( P_f(x) = P_\infty \in \mathbb{R} \) is constant. \( \square \)

Remark 2.2. The proof of Lemma 3 takes advantage of the dimension being one. In higher dimension, the argument above would only show that \( m_f \) is a divergence free vector field, but not necessarily constant. We shall make further use of the one dimensionality of the model when proving pointwise bounds on the NESS.

2.2. Higher moments. Multiplying the steady-state equation

\[
(2.4) \quad v \partial_x f = \alpha M_f + (1 - \alpha) \rho_f \left( \frac{M_{T_1} + M_{T_2}}{2} \right) - f
\]

by \( v^2 \), and integrating in \( v \) yields

\[
\frac{d}{dx} \int_{\mathbb{R}} v^3 f(x, v) \, dv = (1 - \alpha) \left( \frac{T_1 + T_2}{2} \right) (\rho_f(x) - 1),
\]

since by Lemma 3 \( P_f(x) = (T_1 + T_2)/2 \). Next, multiplying (2.4) by \( v^3 \) and integrating yields

\[
\frac{d}{dx} \int_{\mathbb{R}} v^4 f(x, v) \, dv = -\int_{\mathbb{R}} v^3 f(x, v) \, dv.
\]

Combining the last two equations yields

\[
(2.5) \quad -\frac{d^2}{dx^2} \int_{\mathbb{R}} v^4 f(x, v) \, dv = (1 - \alpha) \left( \frac{T_1 + T_2}{2} \right) (\rho_f(x) - 1).
\]

Since the right hand side integrates to zero, we have

\[
\int_{\mathbb{R}} v^4 f(x, v) \, dv - \int_{T \times \mathbb{R}} v^4 f(x, v) \, dv \, dx \, dv = (1 - \alpha) \left( \frac{T_1 + T_2}{2} \right) \int_{T} \psi(x - y)(\rho_f(y) - 1) \, dy
\]

where

\[
\psi(x) = \sum_{k \neq 0} \frac{e^{2\pi i k x}}{4\pi^2 k^2} \quad \text{so that} \quad |\psi(x)| \leq \frac{1}{12}.
\]
It follows that
\begin{equation}
\left| \int_{\mathbb{R}} v^4 f(x, v) \, dv - \int_{\mathbb{R} \times \mathbb{R}} v^4 f(x, v) \, dx \, dv \right| \leq (1 - \alpha) \left( \frac{T_1 + T_2}{12} \right).
\end{equation}

In particular, for \( \alpha \) close to 1, \( \int v^4 f(x, v) \, dv \) is nearly constant; its average is
\begin{equation}
\hat{T} \times R_v^4 \int f(x, v) \, dv = 3 \left[ \alpha \left( \frac{T_1 + T_2}{2} \right)^2 + (1 - \alpha) \frac{T_1^2 + T_2^2}{2} \right],
\end{equation}
and thus the spatial fluctuations in \( \int v^4 f(x, v) \, dv \) are a small fraction of the mean for large temperatures.

**Lemma 4.** Let \( f \) be a solution to \((2.4)\) such that
\begin{equation}
\int_{\mathbb{R} \times \mathbb{R}} (1 + v^2) f(x, v) \, dx \, dv < \infty
\end{equation}
and recall that \( P_\infty := (T_1 + T_2)/2 \). Then \((2.6)\) is valid, and the spatial density \( \rho_f \) satisfies
\begin{equation}
\rho_f(x) \geq \frac{1}{3(2 - \alpha) + \frac{(1 - \alpha)}{6P_\infty}}.
\end{equation}

**Proof.** By the Cauchy-Schwarz inequality and Lemma 3,
\begin{equation}
\forall \ x \in \mathbb{T}, \quad P_\infty = P_f(x) = \int_{\mathbb{R}} v^2 f(x, v) \, dv \leq \left( \int_{\mathbb{R}} v^4 f(x, v) \, dv \right)^{1/2} \rho_f^{1/2}(x),
\end{equation}
so that
\begin{equation}
\rho_f(x) \geq P_\infty^2 \left( \int_{\mathbb{R}} v^4 f(x, v) \, dv \right)^{-1}.
\end{equation}

From \((2.7)\), we have the bounds
\begin{equation}
\left\{ \begin{array}{l}
\int_{\mathbb{R} \times \mathbb{R}} v^4 f(x, v) \, dx \, dv \leq 3(2 - \alpha)P_\infty^2 \\
\sup_{x \in \mathbb{T}} \int_{\mathbb{R}} v^4 f(x, v) \, dv \leq 3(2 - \alpha)P_\infty^2 + \frac{(1 - \alpha)P_\infty}{6}.
\end{array} \right.
\end{equation}
Combining bounds yields the result. \(\square\)

**2.3. Splitting between odd and even parts and a wave-like system.** We split a given steady state \( f \) into even and odd parts \( f = E + O \) with respect to the \( v \) variable. The steady state equation \((1.1)\) can be rewritten as:
\begin{equation}
\left\{ \begin{array}{l}
v \partial_x E = -O \\
v \partial_x O = F_\alpha - E
\end{array} \right.
\end{equation}
where
\begin{equation}
F_\alpha(x, v) := \alpha M_f(x, v) + (1 - \alpha) \rho_f(x) G(v), \quad G := \left( \frac{Mf_1 + Mf_2}{2} \right).
\end{equation}
Combining the two equations in (2.9), we obtain
\[(1 - v^2 \partial_x^2) E = F_\alpha.\]

Note that for each \(v \neq 0\), the operator \((1 - v^2 \partial_x^2)\) is invertible with a bounded kernel. The equation (2.11) conveniently and efficiently expresses the iterated effects of velocity averaging on the steady state, or rather on its even part.

**Lemma 5.** For each density \(\rho\) on \(\mathbb{T}\), there is at most one NESS \(f\) such that \(\rho = \rho_f\).

**Proof.** We have seen that for any NESS, \(P_\infty = (T_1 + T_2)/2\), and then by Lemma 3, \(\rho_f T_f = P_\infty\), so that \(M_f\) is determined by \(\rho\). It follows that \(F_\alpha\) is determined by \(\rho\), and then since (2.11) is uniquely solvable, the uniqueness of \(f\) follows. \(\square\)

The formal solution of (2.12) is
\[(2.12) \quad E = (1 - v^2 \partial_x^2)^{-1} F_\alpha\]

and can be written in terms of an explicit Green’s function. Integrating in \(v\) yields
\[(2.13) \quad \rho_f(x) = \int_{\mathbb{R}} \left((1 - v^2 \partial_x^2)^{-1} F_\alpha\right) \, dv.\]

**Lemma 6.** Let \(T_1, T_2 > 0\) and \(\alpha \in (0, 1)\) and \(P_\infty = (T_1 + T_2)/2\). For any probability density \(\rho = \rho(x)\) on \(\mathbb{T}\), define \(T(x) = P_\infty/\rho(x)\) and
\[(2.14) \quad \begin{cases} M[\rho](x, v) := \frac{\rho(x)}{\sqrt{2\pi T(x)}} e^{-\frac{x^2}{2 T(x)}} = \frac{\rho^{3/2}(x)}{\sqrt{2\pi P_\infty}} e^{-\frac{x^2}{4 P_\infty}} \\ F_\alpha[\rho](x, v) := \alpha M[\rho](x, v) + (1 - \alpha) \rho(x) G(v) \\ \Psi_\alpha[\rho](x) := \int_{\mathbb{R}} \left[(1 - v^2 \partial_x^2)^{-1} F_\alpha[\rho]\right] \, dv \end{cases}
\]

with \(G\) defined as in (2.10). Then for all \(\rho\), \(\Psi_\alpha[\rho]\) is a probability density on \(\mathbb{T}\), and \(\rho\) is the spatial density of some NESS \(f\) if and only if \(\rho = \Psi_\alpha[\rho]\), and in this case the unique such NESS is given in terms of \(\rho\) by (2.15) and (2.16) below.

**Proof.** Let \(\rho\) be any density. Then \(F_\alpha[\rho]\) is a probability density on \(\mathbb{T} \times \mathbb{R}\). Since \((1 + v^2 \partial_x^2)^{-1}\) preserves both integrals and positivity, \((1 + v^2 \partial_x^2)^{-1} F_\alpha[\rho]\) is also a a probability density on \(\mathbb{T} \times \mathbb{R}\), and hence its marginal, \(\Psi_\alpha[\rho]\), is a probability density on \(\mathbb{T}\).

Next, suppose that \(\rho(x) = \rho_f(x)\) for some NESS \(f(x, v)\). By Lemma 3, \(M[\rho](x, v) = M_f(x, v)\), and therefore \(F_\alpha[\rho](x, v)\) is given in terms of \(f\) by (2.10). Then \(E(x, v)\), given by (2.12), is a probability density on \(\mathbb{T} \times \mathbb{R}\). \(E(x, v)\) is the even part of \(f(x, v)\) and finally \(O(x, v)\), the odd part of \(f(x, v)\), is given by the first equation in (2.9). Then by (2.13), \(\rho = \Psi_\alpha[\rho]\).

Finally, suppose that \(\rho(x)\) is a probability density on \(\mathbb{T}\) such that \(\rho = \Psi_\alpha[\rho]\). Define \(F_\alpha[\rho]\) by (2.14), and then define \(E_\alpha[\rho](x, v)\) by
\[(2.15) \quad E_\alpha[\rho] := (1 - v^2 \partial_x^2)^{-1} F_\alpha[\rho],\]
and then define
\begin{equation}
O_\alpha[\rho] = -v \partial_x E_\alpha[\rho] \quad \text{and} \quad f_\rho(x, v) = E_\alpha[\rho](x, y) + O_\alpha[\rho](x, y).
\end{equation}

Then
\begin{align*}
v \partial_x f_\rho(x, v) &= -O_\alpha[\rho] - v^2 \partial_x^2 E_\alpha[\rho] \\
&= -O_\alpha[\rho] - E_\alpha[\rho] + (1 - v^2 \partial_x^2) E_\alpha[\rho] \\
&= F_\alpha[\rho] - f_\rho.
\end{align*}

Lemma 3 applies to this equation, and we conclude that \( \rho f_\rho = P_\infty \). The fixed point equation \( \rho = \Psi_\alpha[\rho] \) implies \( \rho f_\rho = \rho \) and \( \mathcal{M}[\rho] = \mathcal{M}[f_\rho] \). This shows that \( f_\rho \) is an NESS for our equation, and concludes the proof that \( \rho \) is the spatial density of an NESS if and only if it is a fixed point of \( \Psi_\alpha \). \qed

It follows from Lemma 6 that our conjecture would be proved if we could show that the constant density is the only fixed point of \( \Psi_\alpha \) for all \( \alpha \in [0, 1] \). We prove this for sufficiently small \( \alpha \) in the next section.

3. Pointwise bounds on the moments of the NESS

3.1. Preliminaries. We define the standard Fourier series of an integrable function \( r = r(x) \) on the torus \( T = [-1/2, 1/2] \) by
\begin{equation}
\hat{r}(k) := \int_0^1 r(x)e^{-2i\pi k x} \, dx, \quad k \in \mathbb{Z},
\end{equation}
and we recall the inversion formula (when, say, the Fourier modes \( \hat{r}(k) \) for \( k \in \mathbb{Z} \) are absolutely summable)
\begin{equation}
r(x) = \sum_{k \in \mathbb{Z}} \hat{r}(k)e^{2i\pi k x}.
\end{equation}

Define \( \varphi_v(x) \) the fundamental solution to the Laplace equation
\begin{equation}
(1 - (v \partial_x)^2)^{-1} \varphi_v(x) = \delta_0(x)
\end{equation}
on the circle \( T = [-1/2, 1/2] \). This fundamental solution is explicit:
\begin{equation}
\varphi_v(x) = \sum_{m \in \mathbb{Z}} \phi_v(x + m), \quad \phi_v(x) := \frac{1}{2|v|} e^{-\frac{|x|}{|v|}},
\end{equation}
and its formula in Fourier is
\begin{equation}
\hat{\varphi}_v(k) = \int_{-1/2}^{1/2} e^{-2i\pi k x} \varphi_v(x) \, dx = \frac{1}{1 + (2\pi)^2 v^2 k^2}, \quad k \in \mathbb{Z}.
\end{equation}

The following bounds will be useful: \( \varphi_v \in L^p(T) \) for all \( p \in [1, +\infty] \), and
\begin{equation}
\|\varphi_v\|_{L^p(T)} = \left( \frac{1}{p} \right)^\frac{1}{p} \left( \frac{1}{2|v|} \right)^{\frac{p-1}{p}}, \quad p \in [1, +\infty), \quad \|\varphi_v\|_{L^\infty(T)} = \left( \frac{1}{2|v|} \right).
\end{equation}

and it satisfies the lower bound
\begin{equation}
\forall x, y \in T, \quad \varphi_v(x - y) \geq \frac{1}{2|v|} e^{-\frac{1}{|v|}}.
\end{equation}
3.2. Lower bound on the NESS. We have already proved a uniform lower bound on the spatial density of any NESS in Lemma 4. We now prove a stronger result: a uniform lower bound holds for every density in the range of \( \Psi_\alpha \).

**Lemma 7** (Pointwise lower bound). Let \( \alpha \in [0,1] \) and let \( \rho \) be any probability density on \( T \) of the form \( \rho = \Psi_\alpha[\rho_0] \) for a probability density \( \rho_0 \) on \( T \). Then

\[
\forall x \in T, \quad \rho(x) \geq r_\infty := \frac{(1 - \alpha)}{4\sqrt{e}} \int_{1 \leq |v| \leq 2} \left( \frac{M_{\mathcal{T}_1} + M_{\mathcal{T}_2}}{2} \right) \, dv.
\]

**Proof.** Define \( F_\alpha[\rho_0] \) in terms of \( \rho_0 \) using (2.14). Then

\[
\rho(x) = \hat{T} \times R_{x} F_\alpha[\rho_0](y, v) \, dy \, dv.
\]

The operator \( (1 - (v \partial_x)^2)^{-1} \) preserves positivity, and since \( F_\alpha[\rho_0] \geq (1 - \alpha)\rho_0 G \), we obtain

\[
\rho(x) \geq (1 - \alpha) \int \varphi_v(x - y) F_\alpha[\rho_0](y, v) \, dy \, dv
\]

The kernel \( \varphi_v(x - y) \) is bounded below by \( e^{-1/|v|/(2|v|)} \). The function defined by \( t > 0 \mapsto t^{-1}e^{-t^{-1}} \) vanishes as \( t \) approaches zero or infinity, is maximum at \( t = 1 \) and then decreases as \( t \) increases, and has the value \( 1/(2\sqrt{e}) \) at \( t = 2 \). This implies

\[
\varphi_v(x - y) \geq \frac{1}{4\sqrt{e}}
\]

for all \( x, y \) whenever \( 1 \leq |v| \leq 2 \), which concludes the proof. \( \square \)

3.3. Upper bound on the NESS.

**Lemma 8** (Gain of integrability of \( \Psi_\alpha \)). Let \( \alpha \in [0,1] \) and \( r \in [0,1) \), and let \( \rho_0 \in L^{1+r/2}(T) \) a probability density. Then \( \rho := \Psi_\alpha[\rho_0] \in L^{1+r}(T) \), and there are \( A_r, B_r > 0 \) depending only on \( r \) and degenerating as \( r \to 1 \) such that

\[
\int_T \rho^{1+r} \, dx \leq \alpha A_r \int_T \rho_0^{1+r/2} \, dx + B_r.
\]

As a consequence, if \( \rho = \rho_0 \in L^{1+r/2}(T) \) is a fixed point of \( \Psi_\alpha \), there is a constant \( K_{r,\alpha} > 0 \) depending only on \( r \) and \( \alpha \) and monotone increasing in \( \alpha \) such that

\[
\int_T \rho^{1+r} \, dx \leq K_{r,\alpha}.
\]

**Remark 3.1.** Note that the constant \( K_{r,\alpha} \) is independent of \( \|\rho\|_{L^{1+r/2}(T)} \). That is, knowing only that \( \int_T \rho^{1+r/2} \, dx \) is finite, we obtain a universal bound on \( \int_T \rho^{1+r} \, dx \).
Proof. Define again $F_{\alpha}[\rho_0]$ in terms of $\rho_0$ using (2.14) so that

$$
(3.6) \quad \rho(x) = \int_{T \times \mathbb{R}} \varphi_v(x - y) F_{\alpha}[\rho_0](y, v) \, dy \, dv.
$$

Multiply (3.6) by $\rho^r(x)$ and integrate in $x$:

$$
\int_T \rho^{1+r}(x) \, dx = \int_T \rho^r(x) \left( \int_{T \times \mathbb{R}} \varphi_v(x - y) F_{\alpha}[\rho_0](y, v) \, dy \, dv \right) \, dx
$$

$$
= \int_{T \times \mathbb{R}} \left( \int_T \varphi_v(x - y) \rho^r(x) \, dx \right) F_{\alpha}[\rho_0](y, v) \, dy \, dv.
$$

Equation (3.1) implies

$$
\int_T \varphi_v(x - y) \rho_0^r(x) \, dx \leq \|\rho_0\|_{L^1(T)} \|\varphi_v\|_{L^{1+r}(T)} \leq \frac{1}{(2|v|)^r}.
$$

Therefore

$$
\int_T \rho^{1+r}(x) \, dx \leq \int_{T \times \mathbb{R}} |v|^{-r} F_{\alpha}[\rho_0](y, v) \, dy \, dv
$$

$$
\leq \int_{T \times \mathbb{R}} |v|^{-r} \left[ \alpha \mathcal{M}_{\rho_0}(y, v) + (1 - \alpha)\rho_0(y)G(v) \right] \, dy \, dv
$$

$$
\leq \alpha \int_{T \times \mathbb{R}} |v|^{-r} \mathcal{M}_{\rho_0}(y, v) \, dy \, dv + \int_{\mathbb{R}} |v|^{-r} G(v) \, dv.
$$

Now using the definition of $\mathcal{M}_{\rho_0}$ and Lemma 3:

$$
\int_{T \times \mathbb{R}} |v|^{-r} \mathcal{M}_{\rho_0} \, dy \, dv \leq \int_T \left( \int_{\mathbb{R}} |v|^{-r} \frac{\rho_0^{3/2}(y)}{2\pi P_\infty} e^{-\frac{1}{2} \frac{v^2}{P_\infty}} \, dv \right) \, dx
$$

and making the change of variable $w = (\rho_0(y)/P_\infty)^{1/2}v$,

$$
\int_{\mathbb{R}} |v|^{-r} \frac{\rho_0^{3/2}(y)}{2\pi P_\infty} e^{-\frac{1}{2} \frac{v^2}{P_\infty}} \, dv = \frac{\rho_0(y)^{1+r/2}}{2\pi P_\infty} \int_{\mathbb{R}} |w|^{-r} e^{-\frac{|w|^2}{2}} \, dw.
$$

This yields (3.5) with

$$
A_r := \frac{1}{2\pi P_\infty} \int_{\mathbb{R}} |w|^{-r} e^{-\frac{|w|^2}{2}} \, dw \quad \text{and} \quad B_r := \int_{\mathbb{R}} |v|^{-r} G(v) \, dv.
$$

Now suppose that $\rho = \rho_0 \in L^{1+r/2}(T)$ is a fixed point of $\Psi_\alpha$. We have from the inequality and Hölder’s inequality that

$$
\int_T \rho^{1+r} \, dx \leq \alpha A_r \int_T \rho^{1+r/2} \, dx + B_r \leq \alpha A_r \left( \int_T \rho^{1+r} \, dx \right)^{1+r/2} + B_r.
$$

Since the exponent $(1+r/2)/(1+r)$ on the right is less than one, this proves that there is a constant $K_{r,\alpha} > 0$ depending only on $r$ and $\alpha$ such that $\int_T \rho^{1+r} \, dx \leq K_{r,\alpha}$. □
Lemma 9 (Pointwise upper bound). Let $\alpha \in [0,1]$ and let $\rho$ be any probability density on $\mathbb{T}$ that is a fixed point of $\Psi_\alpha$, and such that $\rho \in L^p$ for some $p > 1$. Then $\rho$ satisfies the pointwise upper bound $\rho(x) \leq R_\infty$ where $R_\infty < \infty$ only depends on the total energy, $\alpha$ and on $p$, and is monotone increasing in $\alpha$.

Proof. In the case where $p \in (1,7/4)$, finitely many applications of the previous lemma will yield, for any $q \in (p,2)$, a bound
\[ \int_\mathbb{T} \rho^q(x) \, dx \leq C_{r,\alpha} \]
for some finite constant $C_{r,\alpha}$ depending only on $r$ and $\alpha$. We deduce that for all $p > 1$, the following control holds
\[ \int_\mathbb{T} \rho^{7/4}(x) \, dx \leq C_{r,\alpha}. \]

We return to (3.6) and expand $F_\alpha$ to write
\begin{equation}
\rho(x) = \alpha \int_{\mathbb{T} \times \mathbb{R}} \varphi_v(x-y)M_\rho(x,v) \, dy \, dv + (1-\alpha) \int_{\mathbb{T} \times \mathbb{R}} \varphi_v(x-y)\rho(y)G(v) \, dy \, dv.
\end{equation}

Observe that Lemma 7 implies
\[ M_\rho(x,v) = \frac{\rho^{3/2}(x)}{\sqrt{2\pi P_\infty}} e^{-v^2/2P_\infty} \leq \frac{\rho_0^{3/2}(x)}{\sqrt{2\pi P_\infty}} e^{-r_\infty v^2/2P_\infty}. \]

which yields
\begin{equation}
\int_{\mathbb{T} \times \mathbb{R}} \varphi_v(x-y)M_f(x,v) \, dy \, dv \leq \frac{1}{\sqrt{2\pi P_\infty}} \int_{\mathbb{R}} \exp \left[ -\frac{r_\infty v^2}{2P_\infty} \right] \left( \int_\mathbb{T} \varphi_v(x-y)\rho^{3/2}(y) \, dy \right) \, dv.
\end{equation}

We apply Hölder’s inequality with conjugate exponents $p = 7/6$ and $q = 7$ to obtain
\begin{align*}
\left( \int_\mathbb{T} \varphi_v(x-y)\rho^{3/2}(y) \, dy \right) &\leq \| \varphi_v \|_{L^{7}(\mathbb{T})} \| \rho^{3/2} \|_{L^{7/6}(\mathbb{T})} \\
&\leq \| \varphi_v \|_{L^{7}(\mathbb{T})} \| \rho \|_{L^{7/4}(\mathbb{T})}^{3/2} \\
&\leq |v|^{-6/7} \| \rho \|_{L^{7/4}(\mathbb{T})}^{3/2} \\
&\leq C_{r,\alpha} |v|^{-6/7}
\end{align*}

where in the final step we have the used estimate (3.1) with $p = 7$. Using this in (3.8) and noting that
\[ \int_{\mathbb{R}} \exp \left[ -\frac{r_\infty v^2}{2P_\infty} \right] |v|^{-6/7} \, dv < \infty, \]
we deduce a universal upper bound on $\int_{\mathbb{T} \times \mathbb{R}} \varphi_v(x-y)M_f(x,v) \, dy \, dv$. 

The term \( \int_{\mathbb{T} \times \mathbb{R}} \varphi_v(x-y) \rho(y) G(v) \, dy \, dv \) is bounded using Hölder inequality and (3.1):
\[
\int_{\mathbb{T}} \varphi_v(x-y) \rho(y) \, dy \leq \| \varphi_v \|_{L^{7/3}(\mathbb{T})} \| \rho \|_{L^{7/4}(\mathbb{T})} \leq C_{r, \alpha} |v|^{-4/7}.
\]
The two last inequalities combined with (3.7) imply the pointwise bound on \( \rho \). \( \square \)

4. THE CONTRACTION MAPPING ARGUMENT

4.1. Setting of the argument. Recall the relation \( \rho_f T_f = P_\infty \) and
\[
\begin{align*}
F_\alpha[\rho] &= \alpha \rho M_T \hat{\rho} + (1-\alpha) \rho G, \\
G &= \left( M_{T_1} + M_{T_2} \right) / 2, \\
\Psi_\alpha[\rho] &= \int_{\mathbb{R}} (1-v^2 \partial_x^2)^{-1} F_\alpha[\rho] \, dv.
\end{align*}
\]
The local density of any steady state must be a fixed point of \( \Psi_\alpha \). When \( \alpha = 0 \), the map \( \Psi_0 \) is linear, and a consequence of the spectral analysis of the next Section 5 is that it is strictly contractive in \( H^1 \) or \( L^2 \) norms. To extend it to small positive \( \alpha \), we make use of the \( a\text{-}priori \) bounds proved in the previous Section 3

4.2. The contraction estimate.

**Lemma 10.** Given any \( \varepsilon \in (0,1) \), define
\[
C_\varepsilon := \left\{ \rho \in L^2(\mathbb{T}) : \int_{\mathbb{T}} \rho(x) \, dx = 1 \text{ and } 0 < \varepsilon < \rho < 1/\varepsilon \right\}.
\]
Then \( \Psi_\alpha(C_\varepsilon) \subset L^2(\mathbb{T}) \) for all \( \alpha \in [0,1] \) and there are \( \alpha_\varepsilon, \delta_\varepsilon \in (0,1) \) depending on \( \varepsilon \) such that for all \( \alpha \in [0,\alpha_\varepsilon) \):
\[
\forall \rho_1, \rho_2 \in C_\varepsilon, \quad \| \Psi_\alpha(\rho_1) - \Psi_\alpha(\rho_2) \|_{L^2(\mathbb{T})} \leq (1-\delta_\varepsilon) \| \rho_1 - \rho_2 \|_{L^2(\mathbb{T})}.
\]

**Proof.** Recall that in our normalization the global conserved quantities satisfy
\[
\rho_\infty := \int_{\mathbb{T} \times \mathbb{R}} f(x,v) \, dx \, dv = 1, \quad T_\infty := \frac{1}{\rho_\infty} \int_{\mathbb{T} \times \mathbb{R}} v^2 f(x,v) \, dx \, dv = \frac{T_1 + T_2}{2}.
\]
The fact that \( \Psi_\alpha(C_\varepsilon) \subset L^2(\mathbb{T}) \) is straightforward and we only prove the contraction property. We linearize the map \( \Psi_\alpha \) around a profile \( \bar{\rho} \in C_\varepsilon \) with global mass 1 and global temperature \( T_\infty \). The local temperature is \( T(x) = T_\infty / \bar{\rho}(x) \). We write the fluctuation
\[
\rho = \bar{\rho} + \sigma \quad \text{with} \quad \sigma \in L^2(\mathbb{T}) \text{ and } \int_{\mathbb{T}} \sigma(x) \, dx = 0.
\]
The functional derivative of \( \Psi_\alpha \) is:
\[
D \Psi_\alpha[\bar{\rho}] \cdot \sigma = \alpha \int_{\mathbb{R}} \left[ 1 - v^2 \partial_x^2 \right]^{-1} \left( \frac{3\sigma}{2} M_T \bar{\rho} - \frac{v^2 \bar{\rho} \sigma}{2T_\infty} M_T \right) \, dv \\
+ (1-\alpha) \int_{\mathbb{R}} \left[ 1 - v^2 \partial_x^2 \right]^{-1} \sigma G(v) \, dv.
\]
We estimate by duality for $\sigma_1, \sigma_2 \in L^2(T)$:

$$
\langle \sigma_2, D\Psi_\alpha[\hat{\rho}] \cdot \sigma_1 \rangle = \alpha \int_{T \times \mathbb{R}} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right) \left( \frac{3\sigma_1}{2} M^\infty \rho - \frac{v^2 \hat{\rho} \sigma_1}{2T^\infty} M^\infty \rho \right) \, dx \, dv
+ (1 - \alpha) \int_{T \times \mathbb{R}} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right) G(v) \sigma_1(x) \, dx \, dv
=: D_1 + D_2.
$$

Let us study the first term $D_1$. Using the controls $\hat{\rho} \in C_\varepsilon$ we deduce

$$
\left\lfloor \int_{T} \left( \frac{3}{2} M^\infty \rho + \frac{v^2 \hat{\rho} \sigma_1}{2T^\infty} M^\infty \rho \right)^2 \sigma(x) \, dx \right\rfloor^{1/2} \leq C_\varepsilon e^{-\nu_\varepsilon v^2} \|\sigma\|_{L^2(T)}
$$

for some constant $C_\varepsilon, \nu_\varepsilon > 0$ depending on $\varepsilon$, and therefore

$$
\int_{T \times \mathbb{R}} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right) \left( \frac{3}{2} M^\infty \rho - \frac{v^2 \hat{\rho} \sigma_1}{2T^\infty} M^\infty \rho \right) \sigma_1(x) \, dx \, dv
\leq C_\varepsilon \|\sigma_1\|_{L^2(T)} \int_{\mathbb{R}} \left( \int_{T} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right)^2 \, dx \right)^{1/2} e^{-\nu_\varepsilon v^2} \, dv
\leq C_\varepsilon \|\sigma_1\|_{L^2(T)} \|\sigma_2\|_{L^2(T)}
$$

where we have used that $\|\varphi_v\|_{L^1(T)} = 1$. We conclude that

$$
D_1 \leq \alpha C_\varepsilon \|\sigma_1\|_{L^2(T)} \|\sigma_2\|_{L^2(T)}.
$$

Let us now study the second term $D_2$. We Fourier transform it in $x$:

$$
\int_{T \times \mathbb{R}} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right) G(v) \sigma_1(x) \, dx \, dv = \sum_{k \in \mathbb{Z}} \int_{\mathbb{R}} \frac{\hat{\sigma}_2(k)}{1 + (2\pi)^2 v^2 k^2} G(v) \hat{\sigma}_1(k) \, dv.
$$

The perturbation $\sigma$ has zero mass, hence $\hat{\sigma}(0) = 0$ and

$$
\left| \int_{T \times \mathbb{R}} \left( [1 - v^2 \partial_x^2]^{-1} \sigma_2 \right) G(v) \sigma_1(x) \, dx \, dv \right| \leq \sum_{k \in \mathbb{Z}, \ k \neq 0} \int_{\mathbb{R}} \frac{\hat{\sigma}_2(k)}{1 + (2\pi)^2 v^2 k^2} G(v) \hat{\sigma}_1(k) \, dv
\leq \left( \int_{\mathbb{R}} \frac{G(v)}{1 + (2\pi)^2 v^2} \, dv \right) \sum_{k \in \mathbb{Z}, \ k \neq 0} \hat{\sigma}_1(k) \hat{\sigma}_2(k)
\leq (1 - \delta_G) \|\sigma_1\|_{L^2(T)} \|\sigma_2\|_{L^2(T)}
$$

with $\delta_G \in (0, 1)$, where we have used that

$$
\left( \int_{\mathbb{R}} \frac{G(v)}{1 + (2\pi)^2 v^2} \, dv \right) < \left( \int_{\mathbb{R}} G(v) \, dv \right) = 1.
$$
Therefore the operator $D\Psi_\alpha[\hat{\rho}]$ is bounded from $L^2(T) \to L^2(T)$ with the operator norm bounded by

$$|||D\Psi_\alpha[\hat{\rho}]||| = \sup_{\|\sigma_1\|_{L^2(T)} \leq 1, \|\sigma_2\|_{L^2(T)} \leq 1} \langle \sigma_2, D\Psi_\alpha[\hat{\rho}] \cdot \sigma_1 \rangle \leq \alpha C_\varepsilon + (1 - \alpha)(1 - \delta G).$$

For $\alpha$ small enough we deduce $|||D\Psi_\alpha[\hat{\rho}]||| < (1 - \delta\varepsilon)$ with $\delta\varepsilon \in (0, 1)$. Finally, since the set $C_\varepsilon$ is convex, the mean value theorem gives, for $\rho_1, \rho_2 \in C_\varepsilon$:

$$\|\Psi_\alpha(\rho_1) - \Psi_\alpha(\rho_2)\|_{L^2(T)} = \left\| \int_0^1 D\Psi_\alpha[(1 - t)\rho_1 + t\rho_2] \cdot (\rho_2 - \rho_1) \, dt \right\|_{L^2(T)} \leq (1 - \delta\varepsilon)\|\rho_1 - \rho_2\|_{L^2(T)}$$

which shows the contraction property for the nonlinear map. □

4.3. **Proof of the main Theorem**

We now prove Theorem 1. By Lemma 7 and Lemma 9 there is a $\varepsilon > 0$, depending only on $T_1$ and $T_2$ such that every steady state of (1.1) with $\alpha < 1/2$ belongs to the set $C_\varepsilon$, defined in Lemma 10. Then by Lemma 11 there is an $\alpha_0$ with $0 < \alpha_0 \leq 1/2$ such that $\Psi_\alpha$ is a strictly contractive mapping on this convex set $C_\varepsilon$, and hence there is a unique fixed point in $C_\varepsilon$. Since there is always one spatially uniform steady state, it is the unique steady state.

5. **Perturbative stability of the spatially uniform NESS**

5.1. **Strategy.** In this section we investigate the perturbative stability of the spatially homogeneous NESS

$$f_{\infty, \alpha} := \alpha M_{\infty} + (1 - \alpha)G(v), \quad G(v) := \frac{M_{T_1} + M_{T_2}}{2}, \quad T_\infty := \frac{T_1 + T_2}{2}.$$  

When the reservoirs have different temperatures and are coupled to the system, that is for $\alpha \in (0, 1)$, there is transfer of heat through collisions from the hot reservoir to the cold reservoir, and there is no detailed balance; i.e., time reversal invariance is broken in the steady state. This is reflected in the fact that the linearized operator is a non-self-adjoint operator on $H_\alpha = L^2(f_{\infty, \alpha}^{-1})$ for $\alpha \in (0, 1)$, as we shall see.

Nonetheless, we shall prove that the linearized collision operator still satisfies a microscopic coercivity inequality (see (5.9)), expressing the dissipative nature of the linearized evolution on the orthogonal complement of the null space of the generator. This fact is striking since we do not derive it, through linearization, from a nonlinear entropy principle, which is the usual source of such inequalities. In our non-equilibrium setting, there is no analog of the $H$-Theorem, and therefore we must prove it by direct analysis of linearized collision operator.

Once the microscopic coercivity is proven, we can prove that our system is hypocoercive by a variety of methods, and we briefly describe two of these.
5.2. Linearization around a spatially homogeneous NESS. Consider densities \( f \) that are close to \( f_{\infty, \alpha} \) with fluctuations denoted

\[
\begin{align*}
\begin{cases} 
  h(x, v) := f(x, v) - f_{\infty, \alpha}(v) \\
  \rho(x) = 1 + \sigma(x) \quad \text{with} \quad \sigma(x) := \int_{\mathbb{R}} h(x, v) \, dv, \\
  P(x) = T_{\infty} + \tau(x) \quad \text{with} \quad \tau(x) := \int_{\mathbb{R}} v^2 h(x, v) \, dv,
\end{cases}
\end{align*}
\]

(5.2)

The fluctuations of the local density and local pressure have zero mean:

\[
\int_{\mathbb{T}} \sigma(x) \, dx = 0 \quad \text{and} \quad \int_{\mathbb{T}} \tau(x) \, dx = 0.
\]

Consider the weighted \( L^2 \) Hilbert space of real valued functions defined by the norm

\[
\| h \|_{\mathcal{H}_\alpha}^2 = \int_{\mathbb{T} \times \mathbb{R}} |h(x, v)|^2 \frac{1}{f_{\infty, \alpha}(v)} \, dx \, dv.
\]

We expand \( \mathcal{M}_f - \mathcal{M}_{f_{\infty, \alpha}} \) to first order in terms of \( h, \sigma, \tau \):

\[
\mathcal{M}_f(x, v) - \mathcal{M}_{f_{\infty, \alpha}}(v) = \frac{(1 + \sigma)^{3/2}(x)}{\sqrt{2\pi(T_{\infty} + \tau(x))}} e^{-\frac{v^2(1+\sigma(x))}{2(T_{\infty} + \tau(x))}} - \frac{1}{\sqrt{2\pi T_{\infty}}} e^{-\frac{v^2}{2T_{\infty}}}
\]

\[
\approx \left( \frac{3}{2} - \frac{v^2}{2T_{\infty}} \right) M_{T_{\infty}}(v)\sigma(x) + \left( -\frac{1}{2T_{\infty}} + \frac{v^2}{2T_{\infty}^2} \right) M_{T_{\infty}}(v)\tau(x)
\]

(5.5)

\[
= M_{T_{\infty}}(v)\sigma(x) + \frac{1}{2} \left( \frac{v^2}{T_{\infty}} - 1 \right) M_{T_{\infty}}(v) \left( \frac{1}{T_{\infty}} \tau(x) - \sigma(x) \right).
\]

The fluctuation \( h = f - f_{\infty, \alpha} \) satisfies the equation

\[
\partial_t h + v \partial_x h = \alpha (\mathcal{M}_f - \mathcal{M}_{f_{\infty, \alpha}}) + (1 - \alpha) \sigma G - h.
\]

To first order we obtain the linearized equation

\[
\partial_t h + Sh = L_\alpha h
\]

with the free streaming operator \( S := v \partial_x \) and the linearized collision operator

\[
L_\alpha h(x, v) := \sigma(x) f_{\infty, \alpha}(x, v) + \frac{\alpha}{2} \left( \frac{\tau(x) - \sigma(x)}{T_{\infty}} \right) \left( \frac{v^2}{T_{\infty}} - 1 \right) M_{T_{\infty}}(v) - h(x, v).
\]

Note that both \( L_\alpha \) and \( \mathcal{H}_\alpha \) depend on \( \alpha \) and that \( L_\alpha \) is bounded on \( \mathcal{H}_\alpha \) for all \( \alpha \in [0, 1] \): observe that \( f_{\infty, \alpha} \geq \alpha M_{T_{\infty}} \) and therefore

\[
\| L_\alpha h \|_{\mathcal{H}_\alpha}^2 \lesssim \| \sigma \|_{L^2(\mathbb{T})}^2 + \| \tau \|_{L^2(\mathbb{T})}^2 \left( \int_{\mathbb{R}} \left( \frac{v^2}{T_{\infty}} - 1 \right) \frac{2}{f_{\infty, \alpha}} M_{T_{\infty}}(v) \, dv \right)
\]

\[
\lesssim \| \sigma \|_{L^2(\mathbb{T})}^2 + \left( \int_{\mathbb{R}} \left( \frac{v^2}{T_{\infty}} - 1 \right) \frac{2}{f_{\infty, \alpha}} M_{T_{\infty}}(v) \, dv \right)
\]

\[
\lesssim \| \sigma \|_{L^2(\mathbb{T})}^2 + \frac{1}{T_{\infty}} \| \tau \|_{L^2(\mathbb{T})}^2 \lesssim \max \left\{ 1, \frac{1}{T_{\infty}} \right\} \| h \|_{\mathcal{H}_\alpha}^2.
\]
5.3. **Microscopic coercivity.** We shall now prove that the null space of \( \mathcal{L} \) is the space of functions \( \sigma(x)f_{\infty,\alpha}(v) \) for \( \sigma \in L^2(\mathbb{T}) \) and prove a spectral gap on the orthogonal of this null space.

**Lemma 11.** Let \( \alpha \in [0,1] \) and \( \mathcal{L}_\alpha \) defined as in (5.8). Then for all \( h \in \mathcal{H}_\alpha \),

\[
(5.9) \quad \langle h, \mathcal{L}_\alpha h \rangle_{\mathcal{H}} \leq -\frac{1-\alpha}{2} \int_{\mathbb{T} \times \mathbb{R}} h(x,v) - f_{\infty,\alpha}(v) \int_{\mathbb{R}} h(x,w) \, dw \right|^2 \frac{1}{f_{\infty,\alpha}(v)} \, dx \, dv.
\]

**Remark 5.1.** If we had taken \( \mathcal{H}_\alpha \) to consist if complex valued functions, we would need a real part on the left side of the inequality since \( \mathcal{L}_\alpha \) is not self-adjoint. Note that for \( \alpha \in [0,1) \), the constant \( \lambda_\alpha \) is strictly positive but \( \lambda_\alpha \to 0 \) as \( \alpha \to 1 \). This reflects the fact, see the proof below, that the dissipativity in the energy mode is lost because there is energy conservation in this limit. In fact in this limit case the microscopic coercivity nevertheless holds once accounting for the larger null space of \( \mathcal{L}_1 \). We are not concerned in this case for which the NESS is already known, and we refer to [1] for a study of the microscopic coercivity and hypocoercivity for the equation of this limit case; This could also be deduced from the abstract results in [12].

**Proof.** Let us define the following orthonormal family in \( L^2(f_{\infty,\alpha}^{-1} dv) \):

\[
(5.10) \quad H_0(v) := f_{\infty,\alpha}, \quad H_1(v) := \frac{1}{\sqrt{T_\infty}}vf_{\infty,\alpha}, \quad H_2(v) := c_\alpha \left( \frac{v^2}{T_\infty} - 1 \right) f_{\infty,\alpha},
\]

where \( c_\alpha > 0 \) is the normalizing constant so that \( \|H_2\|_{\mathcal{H}_\alpha} = 1 \) (one can check that \( c_\alpha^{-2} = 3 \left( \alpha + (1-\alpha) \left( 2 - \frac{T_\infty}{T_\infty^2} \right) \right) - 1 \)). We also define the corresponding orthogonal projections \( \Pi_0, \Pi_1, \Pi_2 \) in \( L^2(f_{\infty,\alpha}^{-1} dv) \) (note that they all depend on \( \alpha \)):

\[
(5.11) \quad \begin{cases}
\Pi_0(h)(v) := \left( \int_{\mathbb{R}} h(w) \, dw \right) H_0(v), \\
\Pi_1(h)(v) := \left( \int_{\mathbb{R}} h(w) \frac{w}{\sqrt{T_\infty}} \, dw \right) H_1(v), \\
\Pi_2(h)(v) := \left( \int_{\mathbb{R}} h(w)c_\alpha \left( \frac{w^2}{T_\infty} - 1 \right) \, dw \right) H_2(v) = \left( \frac{\tau}{T_\infty} - \sigma \right) c_\alpha H_2(v).
\end{cases}
\]

Finally we denote \( \Pi^\perp \) the orthogonal projection on \( \{H_0, H_2\}^\perp \) (note that this projection includes \( \Pi_1 \) in its range).

The linearized collision operator \( \mathcal{L}_\alpha \) can be written using this notation a

\[
(5.12) \quad \mathcal{L}_\alpha h = \Pi_0(h) + \frac{\alpha}{2c_\alpha^2} \Pi_2(h) \frac{M_{T_\infty}}{f_{\infty,\alpha}} - h.
\]

We then compute the Dirichlet form

\[
(\mathcal{L}_\alpha h, h)_{\mathcal{H}_\alpha} = ||\Pi_0(h)||_\mathcal{H}_\alpha^2 + \frac{\alpha}{2c_\alpha^2} \left( \Pi_2(h) \frac{M_{T_\infty}}{f_{\infty,\alpha}}, \Pi_2(h) \right)_{\mathcal{H}_\alpha} + \frac{\alpha}{2c_\alpha^2} \left( \Pi_2(h) \frac{M_{T_\infty}}{f_{\infty,\alpha}}, \Pi^\perp(h) \right)_{\mathcal{H}_\alpha} - ||h||_{\mathcal{H}_\alpha}^2.
\]
where we have used that \( H_2M_{T_\infty}f_{\infty,\alpha}^{-1} \) is orthogonal to \( H_0 \) and \( H_1 \) in \( L^2(f_{\infty,\alpha}^{-1}dv) \).

Let us define

\[
U_\alpha(v) := H_2(v)\frac{M_{T_\infty}(v)}{f_{\infty,\alpha}(v)}.
\]

The projection of this function on \( H_0 \) and \( H_1 \) in \( L^2(f_{\infty,\alpha}^{-1}dv) \) is zero and its projection on \( H_2 \) has coefficient

\[
\int_{\mathbb{R}} U_\alpha(v)H_2(v)\frac{1}{f_{\infty,\alpha}(v)} dv = c^2_\alpha \int_{\mathbb{R}} \left( \frac{v^2}{T_\infty} - 1 \right)^2 M_{T_\infty}(v) dv = 2c^2_\alpha.
\]

Its norm satisfies

\[
\int_{\mathbb{R}} U_\alpha(v)^2\frac{1}{f_{\infty,\alpha}(v)} dv = c^2_\alpha \int_{\mathbb{R}} \left( \frac{v^2}{T_\infty} - 1 \right)^2 M_{T_\infty}(v)^2 dv
\leq \frac{c^2_\alpha}{\alpha} \int_{\mathbb{R}} \left( \frac{v^2}{T_\infty} - 1 \right)^2 M_{T_\infty}(v) dv \leq \frac{2c^2_\alpha}{\alpha}
\]

where we have used in the last line \( f_{\infty,\alpha} \geq \alpha M_{T_\infty} \). We then decompose orthogonally \( U_\alpha = \Pi_2(U_\alpha) + \Pi^\perp(U_\alpha) \) and deduce by Pythagoras’ theorem that

\[
\|\Pi^\perp(U_\alpha)\|^2_{L^2(f_{\infty,\alpha}^{-1})} \leq 4c^4_\alpha \frac{1 - \alpha^2}{\alpha^2}.
\]

We deduce on the one hand that

\[
\frac{\alpha}{2c^2_\alpha} \left\langle \Pi_2(h)\frac{M_{T_\infty}}{f_{\infty,\alpha}}, \Pi^\perp(h) \right\rangle_{\mathcal{H}_\alpha} = \frac{\alpha}{2} \int_{\mathbb{T} \times \mathbb{R}} \left( \frac{\tau}{T_\infty} - \sigma \right)^2 U_\alpha(v)H_2(v) dx dv
= \alpha c^2_\alpha \int_{\mathbb{T}} \left( \frac{\tau}{T_\infty} - \sigma \right)^2 dx = \alpha \|\Pi_2(h)\|^2_{\mathcal{H}_\alpha}.
\]

We deduce on the other hand

\[
\frac{\alpha}{2c^2_\alpha} \left\langle \Pi_2(h)\frac{M_{T_\infty}}{f_{\infty,\alpha}}, \Pi^\perp(h) \right\rangle_{\mathcal{H}_\alpha}
= \frac{\alpha}{2c_\alpha} \int_{\mathbb{T} \times \mathbb{R}} \left( \frac{\tau}{T_\infty} - \sigma \right) \Pi^\perp(U_\alpha)\Pi^\perp(h) dx dv
\leq \frac{\alpha}{2c_\alpha} \left\|\Pi^\perp(U_\alpha)\right\|_{L^2(f_{\infty,\alpha}^{-1})} \left( \int_{\mathbb{T}} \left( \frac{\tau}{T_\infty} - \sigma \right)^2 dx \right)^{1/2} \left\|\Pi^\perp(h)\right\|_{\mathcal{H}_\alpha}
\leq \frac{\alpha}{2c^2_\alpha} \left\|\Pi^\perp(U_\alpha)\right\|_{L^2(f_{\infty,\alpha}^{-1})} \|\Pi_2(h)\|_{\mathcal{H}_\alpha} \|\Pi^\perp(h)\|_{\mathcal{H}_\alpha}
\leq (1 - \alpha^2)^{1/2} \|\Pi_2(h)\|_{\mathcal{H}_\alpha} \|\Pi^\perp(h)\|_{\mathcal{H}_\alpha}
\]

where we have used \((5.14)\) in the last line.

We therefore obtain

\[
\langle L_\alpha h, h \rangle_{\mathcal{H}_\alpha} \leq (\alpha - 1)\|\Pi_2(h)\|^2_{\mathcal{H}_\alpha} + (1 - \alpha^2)^{1/2} \|\Pi_2(h)\|_{\mathcal{H}_\alpha} \|\Pi^\perp(h)\|_{\mathcal{H}_\alpha} - \|\Pi^\perp(h)\|^2_{\mathcal{H}_\alpha}.
\]
The quadratic form on the right hand side is negative for \( \alpha \in [0, 1) \) since then 
\[
(1 - \alpha^2) < 4(1 - \alpha)^2.
\] It degenerates at \( \alpha = 1 \). The matrix of the quadratic form 
\[
(V_1, V_2) \in \mathbb{R}^2 \mapsto (\alpha - 1)V_1 - V_2 + (1 - \alpha^2)^{1/2}V_1V_2
\] has characteristic polynomial 
\[
P(X) = X^2 + (2 - \alpha)X + \frac{1}{4}(1 - \alpha)(3 - \alpha)
\] whose roots are \( \alpha/2 + 1 \pm 1/2 \). The greatest eigenvalue is therefore \( (\alpha - 1)/2 \) which concludes the proof. \( \Box \)

5.4. Hypocoercivity. With the microcoercivity at hand we can now readily prove
prove hypocoercivity: That is, we shall prove that for some constant \( C < \infty \) and
some \( \lambda > 0 \), and solution \( h_t \) in \( \mathcal{H}_\alpha \) of our linearized evolution equation satisfies

\[
(5.15) \quad \|h_t\|_{\mathcal{H}_\alpha} \leq Ce^{-\lambda t}\|h_0\|_{\mathcal{H}_\alpha}.
\]
With this in hand, it is a simple matter to prove the nonlinear stability. We discuss
two approaches to proving (5.15) for our model. One approach applies when the
steady state is symmetric in \( v \), as in our case. As noted in [1], whenever this is
the case, there is a natural orthonormal basis such that in this basis the streaming
operator is represented by uncoupled tridiagonal blocks, while the gain term in \( \mathcal{L}_\alpha \)
is represented by uncoupled lower triangular blocks. This structure permits the
extraction of simple, explicit bounds on \( \mu \). Another approach, developed in [9] is
more abstract and not requiring symmetry of the steady state, provides an efficient
route to (5.15). In this section we prove:

**Theorem 12.** The decay estimate (5.15) is valid with the following explicit values
of \( \lambda \) and \( C \):

1. If \( c_2\sqrt{\frac{1-\alpha}{\frac{1}{2}T_\infty}} < \frac{1}{2} \), we may take \( C = 4 \) and \( \lambda = \frac{1-\alpha}{8} \).
2. If \( c_2\sqrt{\frac{1-\alpha}{\frac{1}{2}T_\infty}} \geq \frac{1}{2} \), we may take \( C = 4 \) and \( \lambda = \frac{\sqrt{T_\infty}}{8} \).

This is also true, with the same constants, if we replace \( \mathcal{H}_\alpha \) by \( \mathcal{H}_\alpha^1 \), the latter
Hilbert space being defined in (1.5).

**Remark 5.2.** This result is stronger than Theorem 2 in that it provides explicit
bounds on the exponential rate of convergence. The reason for the \( \sqrt{T_\infty} \)
dependence of \( \lambda \) for small \( T_\infty \) is that hypocoercivity depends on the effects of the streaming
operator \( v\partial_v \) to “mix” the conserved mass mode into the dissipated modes, and the
tridiagonal representation of the streaming operator given in (5.17) shows that its
mixing effects are proportional to \( \sqrt{T_\infty} \). When \( T_\infty \) is large, there is rapid mixing,
but this can do only so much good: The dissipativity of the non-conserved modes
as estimated in Lemma (11) is bounded independent of \( T_\infty \), but the mixing only
shares this dissipativity around, it cannot improve the dissipativity no matter how
fast it runs.

It is also worth noting that we have simple bounds on \( c_\alpha \): By the arithmetic-
geometric mean inequality, \( 0 \leq \frac{T_\infty}{T_\infty^2} \leq 1 \), and hence \( 2 \leq c_\alpha^2 \leq 3(2 - \alpha) - 1 \).
The proof of Theorem 5.2 is quite short once one has computed matrix representations of $v\partial_x$ and $\mathcal{L}_\alpha$ with respect to a basis that we now introduce: The basis is \( \{ e^{ikx}g_m(v) \}_{m \geq 0, k \in \mathbb{Z}} \) where \( \{ g_m(v) \}_{m \geq 0} \) is the sequence one gets by applying this Gram-Schmidt orthonormalization procedure to the sequence of functions $v^mf_{\infty,\alpha}(v)$ for $m \geq 0$. We write these in the form $g_m(v) = H_m(v)f_{\infty,\alpha}(v)$ where $H_m$ is a polynomial of degree $m$. For $h \in \mathcal{H}_\alpha$, we write

\begin{equation}
(5.16) \quad h = \sum_{m \geq 0, k \in \mathbb{Z}} ik\hat{h}_m(k)e^{ikx}g_m(v).
\end{equation}

The action of the free streaming operator $\mathcal{S} := v\partial_x$ then is

\begin{equation}
\mathcal{S}h(x, v) = \sum_{m \geq 0, k \in \mathbb{Z}} ik\hat{h}_m(k)e^{ikx}vg_m(v).
\end{equation}

It is a simple consequence of the fact that $f_{\infty,\alpha}$ is even in $v$ that for each $m \geq 1$, $vg_m$ is a linear combination of $g_{m-1}$ and $g_{m+1}$; see [1]. Since the operation of multiplication by $v$ is self-adjoint, it follows that there exist numbers \( \{ a_n \}_{m \geq 1} \) such that

\begin{equation}
vg_m(v) = \sqrt{T_\infty}a_{m+1}g_{m+1}(v) + \sqrt{T_\infty}a_{m+1}g_{m-1}(v)
\end{equation}

with the convention $g_{-1} := 0$.

For $m = 0$, by (5.10), $vg_0(v) = vf_{\infty,\alpha} = \sqrt{T_\infty}g_1$, and hence $a_1 = 1$. Likewise,

\begin{equation}
vg_1(v) = \frac{1}{\sqrt{T_\infty}}v^2f_{\infty,\alpha} = \sqrt{T_\infty} \left( \frac{v^2}{T_\infty} - 1 \right) f_{\infty,\alpha} + \sqrt{T_\infty}g_0(v)
\end{equation}

and hence $a_2 = c^{-1}_\alpha$. One can work out $a_m$ for or higher values of $m$, but these are not needed here.

Let $\hat{h}(k)$ denote the element of $\ell_2$ whose $m$th component is $\hat{h}_m(k)$. Then the corresponding vector of coefficients for $\mathcal{S}h$ is given by $ik\hat{S}\hat{h}(k)$ where $\hat{S}$ is the tridiagonal matrix

\begin{equation}
(5.17) \quad \hat{S} = \sqrt{T_\infty} \begin{pmatrix}
0 & a_1 & 0 & 0 & \cdots \\
a_1 & 0 & a_2 & 0 & \cdots \\
0 & a_2 & 0 & a_3 & \cdots \\
0 & 0 & a_3 & 0 & \cdots \\
\vdots & \vdots & \vdots & \vdots & \ddots
\end{pmatrix} = \sqrt{T_\infty} \begin{pmatrix}
0 & 1 & 0 & 0 & \cdots \\
1 & 0 & c^{-1}_\alpha & 0 & \cdots \\
0 & c^{-1}_\alpha & 0 & a_3 & \cdots \\
0 & 0 & a_3 & 0 & \cdots \\
\vdots & \vdots & \vdots & \vdots & \ddots
\end{pmatrix}.
\end{equation}

We now turn to the gain term in the linearized collision operator. The linearized collision operator does not act on the spatial variable. By projecting (5.12) in our basis we get

\begin{equation}
(5.18) \quad \mathcal{L}_\alpha h(k, v) = \hat{h}_0(k)g_0(v) + \frac{\alpha}{2c^2_\alpha} \hat{h}_2(k)H_2(v)MT_\infty(v) - h(k, v).
\end{equation}

For each $k$, the action of $\mathcal{L}_\alpha$ in the \( \{ g_m \}_{m \geq 0} \) basis is given by

\begin{equation}
\mathcal{L}_\alpha h_m(k) = (L_\alpha \hat{h}(k))_m
\end{equation}
where \( L_\alpha + I_d \) is the matrix whose first column is the unit vector \((1, 0, 0, \ldots)\), whose third column is the vector \((0, 0, \alpha, b_3, b_4, \ldots)\), with all other columns being zero, and with

\[
b_m := \frac{\alpha}{2c^2} \hat{h}_2(k) \int H_2(v) H_m(v) M_{\infty}(v) dv ,
\]

so that, in particular, \( b_0 = b_1 = 0 \) and \( b_2 = \alpha \).

Therefore, we may rewrite our linearized equation as the decoupled system of equations

\[
\partial_t \hat{h}(k) = (L_\alpha - ik S) \hat{h}(k) ,
\]

for each \( k \in \mathbb{Z} \). For \( k = 0 \) we simply have from (5.18) \( \partial_t \hat{h}(0) = -\hat{h}(0) \) since \( \hat{h}(0)_0 = \hat{h}(0)_2 = 0 \).

For each \( k \neq 0 \), define \( C_k = -(L_\alpha - ik S) \). We seek a positive definite matrix \( P_k \) such that for some fixed \( \lambda > 0 \),

\[
\forall k \in \mathbb{Z}, \quad C_k^* P_k + P_k C_k \geq 2\lambda P_k.
\]

Then if we define the Lyapunov function \( e(h) \) by

\[
e(h) = \sum_{k \in \mathbb{Z}} \omega_k \left( \hat{h}(k), P_k \hat{h}(k) \right)_{\ell_2},
\]

for any sequence \( \{\omega_k\}_{k \in \mathbb{Z}} \) of positive numbers we have that for any solution of our linearized equation with initial data \( h_0 \) with \( e(h) < \infty \),

\[
\frac{d}{dt} e(h(t)) \leq -2\lambda e(h(t)) \quad \Rightarrow \quad e(h(t)) \leq e^{-2t\lambda} e(h_0).
\]

We will construct the matrices \( P_k \) so that for some \( C > 0 \)

\[
\forall k \in \mathbb{Z}\setminus\{0\}, \quad K I \leq P_k < \frac{1}{K} I.
\]

This implies that the function \( e(h) \) is equivalent to the norm on \( \mathcal{H}_\alpha \) if we take each \( \omega_k = 1 \). We then conclude that (5.15) is valid with \( C = \frac{1}{K} \) and the value of \( \lambda \) appearing in (5.20). By making other choices for \( \omega_k \), we obtain decay in various Sobolev type norms. For instance, taking \( \omega_k = 1 + k^2 \), we would obtain decay in the norm defined in (1.15). With the matrix representation computed, and Lemma 11 at our disposal, we are ready to prove Theorem 12.

**Proof of Theorem 12.** Because \( L_\alpha \) is lower triangular with positive diagonal entries that are uniformly bounded away from zero except of course for the zero in the upper left, and because \( S \) is tridiagonal, a simple prescription from [1] provides \( P_k \). For a parameter \( c \in (0, 1) \) to be chosen later, define \( P_k(a) \) by entering

\[
\begin{pmatrix}
1 & -ic/k \\
-ic/k & 1
\end{pmatrix}
\]

as its upper-left \( 2 \times 2 \) block, with all other entries being those of the identity. The eigenvalues of \( P_k(a) \) are \( 1, 1 + c/k \) and \( 1 - c/k \), and hence (5.22) is satisfied with \( K = 1 - c \).

Then simple computations show that

\[
C_k^* P_k(a) + P_k(a) C_k = (2I - L - L^T) + R
\]
where $R$ is the matrix whose upper $3 \times 3$ block is
\begin{equation}
\begin{pmatrix}
2c\sqrt{T_\infty} & -ic/k & c\sqrt{T_\infty}c^{-1}_\alpha \\
ic/k & 0 & 0 \\
c\sqrt{T_\infty}c^{-1}_\alpha & 0 & 0
\end{pmatrix}
\end{equation}
and whose remaining entries are all zero. Hence for any (real) $h = (h_0, h_1, h_2, \ldots) \in l^2$,
\[
(h, Rh)_{l^2} = 2c\sqrt{T_\infty}h_0^2 + 2c\sqrt{T_\infty}c^{-1}_\alpha h_0 h_2 \geq c\sqrt{T_\infty}(h_0^2 - c^{-2}_\alpha h_2^2).
\]
By Lemma (11),
\[
(h, (2I - L - L^T)h)_{l^2} \geq (1 - \alpha) \sum_{m=1}^{\infty} h_m^2.
\]
Combining these estimates with $I \geq (1 - c)P_k$, (5.20) then holds with $2\lambda = \sqrt{T_\infty} \min \left\{ c, \frac{1 - \alpha}{\sqrt{T_\infty}} - cc^{-2}_\alpha \right\} (1 - c)$.

We now choose $c$ so that $c = \min \left\{ \frac{1}{2}, c^{-2}_\alpha \frac{1 - \alpha}{2\sqrt{T_\infty}} \right\}$. If $c^{-2}_\alpha \frac{1 - \alpha}{2\sqrt{T_\infty}} < \frac{1}{2}$, then $\lambda \geq \frac{1 - \alpha}{8}$. If $c^{-2}_\alpha \frac{1 - \alpha}{2\sqrt{T_\infty}} \geq \frac{1}{2}$, then $\lambda \geq \frac{\sqrt{T_\infty}}{8}$.

We now explain another route to (5.15) which is relies on general abstract results obtained in [9]:

**Theorem 13** (Abstract result of hypocoercivity from [9]). Consider a Hilbert space $H$ and two closed unbounded operators $S$ and $L$ such that:

(H1) **Microscopic coercivity:** There exists an orthogonal projection $\Pi_0$ such that $L\Pi_0 = \Pi_0L = 0$ and there is $\lambda_m > 0$ such that
\[
\forall h \in \text{Domain}(L), \quad -\langle Lh, h \rangle \geq \lambda_m \| (I - \Pi_0)h \|^2.
\]

(H2) **Macroscopic coercivity:** The operator $S$ satisfies $S^* = -S$ (skew symmetry) and there exists $\lambda_M > 0$ such that
\[
\| S\Pi_0 h \|^2 \geq \lambda_M \| \Pi_0 h \|^2 \quad \text{for all } h \in H \text{ such that } \Pi_0(0) \in \text{Domain}(S).
\]

(H3) **Consistency:** $\Pi_0\Pi_0 = 0$.

(H4) **Auxiliary operator:** Define $A := (1 + (\Pi_0)^*S\Pi_0)^{-1}(\Pi_0)^*$ and assume that $AS(1 - \Pi_0)$ and $AL$ are bounded with a constant $C_M > 0$ such that
\[
\forall h \in H, \quad \| AS(1 - \Pi_0)h \| + \| ALh \| \leq C_M \| (1 - \Pi_0)h \|.
\]

Then there exist positive constants $\lambda > 0$ and $C > 0$, which are explicitly computable in terms of $\lambda_m$, $\lambda_M$, and $C_M$, such that
\[
\forall h \in H, \forall t \geq 0, \quad \| e^{t(L-S)}h \| \leq Ce^{-\lambda t} \| h \|.
\]

We give a short proof of here for the convenience of the reader.
Proof of Theorem 13. Define the modified norm
\[ \mathcal{H}[h] := \frac{1}{2} \|h\|^2 + \varepsilon \langle Ah, h \rangle \]
where \( \varepsilon > 0 \) will be chosen small enough below. Given \( h_t := e^{t(L-S)}h \), we compute
\[
\frac{d}{dt} \mathcal{H}[h_t] = \langle Lh_t, h_t \rangle - \varepsilon \langle AS \Pi_0 h_t, h_t \rangle - \varepsilon \langle AS(1 - \Pi_0)h_t, h_t \rangle + \varepsilon \langle SLh_t, h_t \rangle + \varepsilon \langle ALh_t, h_t \rangle
\]
\[= - \mathcal{D}[h_t]. \]
We have used here that \( L^*A = 0 \) which follows from \( A = \Pi_0 A \) and \( \Pi_0 L = 0 \) in (H1). By (H1), (H2), and by \( AS \Pi_0 = (1 + (S \Pi_0)^* (\Pi_0))^{-1}(S \Pi_0)^*(S \Pi_0) \), the sum of the first two terms in \( \mathcal{D}[h_t] \) is coercive:
\[ -\langle Lh_t, h_t \rangle + \varepsilon \langle AS \Pi_0 h_t, h_t \rangle \geq \min \left\{ \lambda_m, \frac{\varepsilon \lambda_M}{1 + \lambda_M} \right\} \|h_t\|^2. \]

Let us prove that the operators \( A \) and \( S \) are bounded:
\[
(5.25) \quad \forall h \in \mathcal{H}, \quad \|Ah\| \leq \frac{1}{2} \|(1 - \Pi_0)h\| \quad \text{and} \quad \|Sh\| \leq \|(1 - \Pi_0)h\|. \]

The equation \( Ah = g = \Pi_0 g \) (remember that \( A = \Pi_0 A \)) is equivalent to
\[ (S \Pi_0)^* h = g + (S \Pi_0)^*(S \Pi_0) g. \]
Taking the scalar product of the above equality with \( g \) and using (H3), we get
\[ \|g\|^2 + \|S \Pi_0 g\|^2 = \langle h, S \Pi_0 g \rangle = \langle (1 - \Pi_0)h, S \Pi_0 g \rangle \leq \|(1 - \Pi_0)h\||S \Pi_0 g\|^2 \leq \frac{1}{4} \|(1 - \Pi_0)h\|^2 + \|S \Pi_0 g\|^2, \]
which completes the proof of (5.25).

The first inequality in (5.26) implies that \( \mathcal{H}[h] \) is equivalent to \( \|h\|^2 \):
\[
(5.26) \quad \frac{1}{2}(1 - \varepsilon)\|h\|^2 \leq \mathcal{H}[h] \leq \frac{1}{2}(1 + \varepsilon)\|h\|^2. \]

The second inequality in (5.25) and (H1)-(H2)-(H3)-(H4) imply
\[
\mathcal{D}[f] \geq \lambda_m \|(1 - \Pi_0)h_t\|^2 + \frac{\varepsilon \lambda_M}{1 + \lambda_M} \|\Pi_0 h_t\|^2 - \varepsilon(1 + C_M)\|(1 - \Pi_0)h_t\|\|h_t\|
\]
\[
\geq \left[ \lambda_m - \varepsilon(1 + C_M) \frac{1}{2\delta} \right] \|(1 - \Pi_0)h_t\|^2 + \varepsilon \left[ \frac{\lambda_M}{1 + \lambda_M} - (1 + C_M)\frac{\delta}{2} \right] \|\Pi_0 h_t\|^2
\]
for an arbitrary \( \delta > 0 \). By choosing first \( \delta \) and then \( \varepsilon \) small enough, a positive constant \( \kappa \) can be found, such that \( \mathcal{D}[h_t] \geq \kappa \|h_t\|^2 \). Using (5.26), this implies
\[
\frac{d}{dt} \mathcal{H}[h_t] \leq -\frac{2\kappa}{1 + \varepsilon} \mathcal{H}[h_t],
\]
completing the proof with \( \lambda = \kappa/(1 + \varepsilon) \) and \( C = \sqrt{1 + \varepsilon}/\sqrt{1 - \varepsilon} \). \( \square \)
Second proof of Theorem 2. We consider \( \alpha \in [0, 1) \) (for \( \alpha = 1 \) the result is already known from [1]). We apply the previous Theorem 13 with \( \mathcal{H} \) being the subspace of \( \mathcal{H}_\alpha = L^2(f_{\infty,0}^{-1} dv) \) consisting of functions that satisfy the zero global mass condition \( \int_{\mathbb{R} \times \mathbb{R}} h(x, v) \, dx \, dv = 0 \). We take \( \mathcal{S} = \nu \partial_x \) and \( \mathcal{L} = \mathcal{L}_\alpha \), the linearized operator defined in (5.12), and \( \Pi_0(h) = (\int_{\mathbb{R}} h \, dv)f_{\infty,0} \) defined in (5.11). Then (H1) is proved in Lemma 11 and (H2) follows from

\[
\int_{\mathbb{R} \times \mathbb{R}} (\nu \partial_x \sigma f_{\infty,0})^2 \, dx \, dv = T_\infty \| \partial_x \sigma \|_{L^2(\mathbb{T})}^2 \geq T_\infty \| \sigma \|_{L^2(\mathbb{T})}
\]

where we have used the Poincaré inequality in the unit torus, and the zero global mass condition. (H3) follows from \( \int_{\mathbb{R}} \nu f_{\infty,0}(v) \, dv = 0 \). Finally, to prove (H4), we first figure out some explicit formula for \( \mathcal{A} \):

\[
\begin{align*}
\mathcal{A} h &= - \left[ (1 - T_\infty \partial_x^2)^{-1} \partial_x \right] f_{\infty,0}, \quad j(x) := \int_{\mathbb{R}} v h(x, v) \, dv \\
\mathcal{A} \mathcal{S} h &= - \left[ (1 - T_\infty \partial_x^2)^{-1} \partial_x^2 \tau \right] f_{\infty,0}, \quad \tau(x) := \int_{\mathbb{R}} v^2 h(x, v) \, dv.
\end{align*}
\]

Since \( (1 - T_\infty \partial_x^2)^{-1} \partial_x \) and \( (1 - T_\infty \partial_x^2)^{-1} \partial_x^2 \) are bounded operators on \( L^2(\mathbb{T}) \), we deduce that \( \mathcal{A} \) and \( \mathcal{A} \mathcal{S} \) are bounded, and (H4) follows since \( \mathcal{L}_\alpha \) is bounded and \( \mathcal{L}_\alpha \Pi_0 = 0 \). This concludes the proof of hypocoercivity for our linearized operator \( \mathcal{L}_\alpha - \mathcal{S} \).

5.5. Nonlinear stability. We close this section by proving nonlinear stability. Let \( \alpha \in [0, 1) \) and consider a probability density \( f_{\infty} \in \mathcal{H}_\alpha \) and define \( h_{\infty} := f_{\infty} - f_{\infty,0} \in \mathcal{H}_\alpha \). This fluctuation has zero global mass by definition. We define the solution through

\[
h_t = e^{(\mathcal{L}_\alpha - \mathcal{S})} h_{\infty} + \int_0^t e^{(t-s)(\mathcal{L}_\alpha - \mathcal{S})} \mathcal{R}[h_s] \, ds
\]

with the nonlinear remainder term defined by

\[
\mathcal{R}[h] := \mathcal{M}_f - \mathcal{M}_{f_{\infty,0}} - \left[ M_{\infty} \sigma + \frac{1}{2} \left( \frac{\nu^2}{T_\infty} - 1 \right) M_{\infty}(v) \left( \frac{1}{T_\infty} \tau - \sigma \right) \right]
\]

with \( \sigma \) and \( \tau \) defined in terms of \( h \) as before. Taylor expansions and straightforward calculations, using the multiplicative property of the \( H^1(\mathbb{T}) \) Sobolev norm, show that

\[
\| \mathcal{R}[h] \|_{\mathcal{H}^1_\alpha} \lesssim \| \sigma \|^2_{H^1(\mathbb{T})} + \| \tau \|^2_{H^1(\mathbb{T})} \lesssim \| h \|^2_{\mathcal{H}^2_\alpha}.
\]

(For more detail in a closely related argument, see [1].) We deduce the a priori estimate

\[
\| h_t \|_{\mathcal{H}_\alpha} \leq C e^{-\lambda t} \| h_{\infty} \|_{\mathcal{H}_\alpha} + C' \int_0^t e^{-\lambda(t-s)} \| h_s \|^2_{\mathcal{H}_\alpha} \, ds
\]

for some constants \( C, C', \lambda > 0 \), and one can use the values of \( C \) and \( \lambda \) provided by Theorem 12. By a standard argument, this shows global existence and exponential decay at rate \( \lambda \) when \( \| h_{\infty} \|_{\mathcal{H}_\alpha} \) is small enough. \qed
Acknowledgements. E. C. acknowledges partial support from NSF grant DMS–174625. J. L. acknowledges support from the AFOS under the award number FA-9500-16-0037 C. M. acknowledges partial support from the ERC grant MAFRAN. He also thanks Rutgers University and the IAS for their invitations, during which this work was started.

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