RESCALED OBJECTIVE SOLUTIONS OF FOKKER–PLANCK AND
BOLTZMANN EQUATIONS

KARSTEN MATTHIES† AND FLORIAN THEIL‡

Abstract. We study the long-time behavior of symmetric solutions of the nonlinear Boltzmann equation and a closely related nonlinear Fokker–Planck equation. If the symmetry of the solutions corresponds to shear flows, the existence of stationary solutions can be ruled out because the energy is not conserved. After anisotropic rescaling, both equations conserve the energy. We show that the rescaled Boltzmann equation does not admit stationary densities of Maxwellian type (exponentially decaying). For the rescaled Fokker–Planck equation we demonstrate that all solutions converge to a Maxwellian in the long-time limit, however, the convergence rate is only algebraic, not exponential.

Key words. objective solution, Boltzmann equation, Fokker–Planck, hypocoercivity

AMS subject classifications. 35Q20, 35Q84, 35K65, 37L20, 70K20

DOI. 10.1137/18M1202335

1. Introduction. Symmetric solutions play a very important role in materials sciences. The reason is that the fundamental laws of physics exhibit many symmetries such as translation and rotation invariance; those symmetries lead to the existence of time-dependent solutions that are invariant under the action of a symmetry group.

The term “objective solution” has been coined by Dumitrica and James in [17] for the case where the symmetry group is a subgroup of the Euclidean symmetry group motivated by molecular dynamics simulations and other engineering applications. We will study objective solutions in the case where the symmetries consist of translations. For the purpose of this paper we say that for a given matrix $S \in \mathbb{R}^{m \times n}$ a function $f : \mathbb{R}^n \to \mathbb{R}$ is $S$-objective if $f(\xi + \eta) = f(\xi)$ for all $\eta \in \ker S$ or, equivalently, if $f(\xi) = g(S \xi)$ for some $g : \mathbb{R}^m \to \mathbb{R}$. We will be mostly interested in the kinetic setting where $\xi = (z, w)$, $z$ being the position and $w$ the velocity. It is important to realize that translation invariance implies that the configuration space is unbounded, therefore extensive thermodynamic quantities such as energy are automatically infinite. Moreover as we are dealing with open systems, it is not necessarily the case that local energy densities are conserved even if the equations of motion are conservative.

The properties of the symmetric solutions depend strongly on the choice of $S$; we analyze here one interesting $S$ which leads to a nonconservative system, but ideas will be also relevant for other $S$. If $n = 2d$, $\text{Id} \in \mathbb{R}^{d \times d}$ is the identity matrix, and $S = (\text{Id}, 0) \in \mathbb{R}^{d \times 2d}$ one obtains solutions that are independent of $\xi$ and the choice $S = (\text{Id}, \pm \text{Id})$ yields expanding and contracting flows where $w = \mp z$. We will study Couette flows/shear flows where

$$S = (-\mu \alpha \otimes \beta, \text{Id})$$

with $\mu \in \mathbb{R}$ being the shear parameter, $\alpha, \beta \in \mathbb{R}^d$ being orthonormal. To see that $S$
corresponds to shear flows observe that

\[ \ker S = \text{span}\{ (\alpha, 0), (\beta, \mu \alpha) \} \]

so that

\[ f(z + x \alpha + y \beta, w + \mu y \alpha) = f(z, w). \]

One of the key obstacles to studying the long-time behavior is the fact that stationary solutions do not exist as the energy density of symmetric solutions increases with time. A popular approach to overcome the problem of energy growth is to consider rescaled objective solutions [20, 10, 11, 17] and, in particular, [23]. We revisit the concept of rescaled objective solutions for the Boltzmann equation and a Fokker–Planck equation with similar properties. In contrast to much of the earlier work, our results are based on the notion of anisotropically rescaled solutions, the nonautonomous anisotropic coordinate change will fix the second moment tensor. We analyze the corresponding rescaled—now nonautonomous—equations and obtain the following results for the nonlinear Fokker–Planck equation (A) and the Boltzmann equation with hard sphere collisions (B).

(A) Characterization of stationary solutions and sharp estimates of the convergence rate (Theorem 5). The convergence rate is algebraic.

(B) Characterization of the collision invariants and a rigorous proof that stationary solutions are not Maxwellian (Theorem 13).

The main difference between the nonlinear Fokker–Planck equation and the Boltzmann equation is that the former has a purely local dissipation term whereas the Boltzmann equation involves a nonlocal and nonlinear collision operator. As a result we can obtain much more detailed information about the long-term behavior of rescaled objective solutions of the Fokker–Planck equation than the Boltzmann equation. In the conservative case it is well known that the Maxwellian is the unique stationary solution of the Fokker–Planck equation and the Boltzmann equation. Moreover solutions of the linear Fokker–Planck equation and the nonlinear, homogeneous Boltzmann equation converge to the equilibrium at an exponential rate; cf. [12] and [29]. For the inhomogeneous Boltzmann equation the problem of establishing exponential convergence to the equilibrium is closely linked to Cercignani’s conjecture; an overview can be found in [14].

The behavior of the rescaled objective solutions is quite different. In the case of the Fokker–Planck equation the equilibrium after the anisotropic scaling is still a Maxwellian, but the rate of convergence is only algebraic. While it is not known whether the rescaled Boltzmann equation for hard spheres admits stationary solutions our results imply that even if one exists it is not of exponential type. In particular, Maxwellians are not equilibria. We point out that existence of renormalized stationary solutions of the Boltzmann equation with Maxwellian interaction has been established in [23].

The main method to analyze the long-term behavior of the Fokker–Planck equation is an adaption of hypocoercivity in a nonautonomous setting. Convergence to equilibria in degenerate dissipative equations preserving mass has attracted major interest starting with the use of logarithmic Sobolev inequalities, entropies, and other functional analytic tools [30, 27]. These methods could be applied to Fokker–Planck equations [1, 7] as well as some Boltzmann equations [4, 13]. A general abstract approach for evolution equations consisting of a (possibly) degenerate dissipative part and some conservative part was introduced by Villani with his concept of hypocoercivity [33]; see also [15]. This method has successfully been adapted in many contexts.

Our methodological contribution is an adaption to nonautonomous nonlinear equations by combining the abstract hypocoercivity result for a limiting problem in a Duhamel formula with a priori estimates for higher derivatives of the full equation. These a priori estimates are indeed obtained using a calculus inspired by hypocoercivity. A crucial ingredient is the detailed asymptotic analysis of the anisotropic rescaling, which can be obtained from closed ordinary differential equations for the second moments of the rescaled Fokker–Planck solutions. Indeed, higher order moment equations are used to derive lower algebraic estimates in the convergence rate for typical initial data. The lack of detailed knowledge about the second moments implies that we have a less explicit control of the anisotropic rescaling in the case of objective solutions to the Boltzmann equation, such that the characterization of a limit distribution and their convergence rates is beyond the scope of this paper.

The rest of the paper is organized as follows. In section 2 we collect some fundamental properties of objective functions. The results for the Fokker–Planck situation are given and proved in section 3. The corresponding analysis for the Boltzmann equation is in section 4. We give a short summary and conclusion in section 5. The proofs of some technical results not relevant for the main argument are postponed to the appendix.

2. Objective functions.

**Definition 1.** Let $S \in \mathbb{R}^{l \times n}$ a matrix. A function $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ is called $S$-objective if $f(\xi + \eta) = f(\xi)$ for all $\eta \in \ker S$.

A classical result for functions which are invariant under the action of a symmetry group is the Hilbert–Weyl theorem which states that the ring of invariant polynomials has a basis; cf., e.g., [21]. We require a closely related result for measurable functions.

**Proposition 2.** Let $S \in \mathbb{R}^{l \times n}$ be a matrix. Let $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ be a measurable function. The following are equivalent:

1. $f$ is $S$-objective.
2. $\nabla \cdot (fT) = 0$ if $T \in \mathbb{R}^{n \times l}$ has the property that $\text{range } T = \ker S$.
3. There exists a measurable function $g : \text{range}(S) \to \mathbb{R}$ such that $f(\xi) = g(S\xi)$.

The proof is standard, we include it for the convenience of the reader.

**Proof.**

1 implies 2: It suffices to show that $\int f \nabla \varphi \cdot \eta \, d\xi = 0$ for each $\eta \in \ker S$ and each smooth and compactly supported test function $\varphi$. As $f$ is $S$-objective one finds that

$$0 = \lim_{h \to 0} \frac{1}{h} \int (f(\xi + h\eta) - f(\xi)) \varphi(\xi) \, d\xi$$

$$= \lim_{h \to 0} \frac{1}{h} \int (\varphi(\xi - h\eta) - \varphi(\xi)) f(\xi) \, d\xi$$

$$= -\int (\nabla \varphi \cdot \eta) f \, d\xi,$$

which is the claim.
2 implies 1: As range $T = \text{ker} S$ there exists $a \in \mathbb{R}^l$ such that $\eta = Ta$. Then
\[
f(\xi + \eta) - f(\xi) = \int_0^1 \frac{d}{ds} f(\xi + s\eta) \, ds = \int_0^1 \nabla f(\xi + s\eta) \cdot \eta \, ds = \int_0^1 \nabla \cdot (f(\xi + s\eta) T) a \, ds = 0.
\]

1 implies 3: Define the operator $\tilde{S} : \text{range}(S^*) \to \text{range}(S)$ by $\tilde{S} = S|_{\text{range}(S^*)}$. Observe that $\tilde{S}$ is invertible and define $g(\xi) = f(\tilde{S}^{-1} \xi)$.

3 implies 1: If $\eta \in \text{ker} S$, then $f(\xi + \eta) = g(S(\xi + \eta)) = g(S\xi) = f(\xi)$.

We are interested in a shear flow setting where $\alpha, \beta \in \mathbb{R}^d$ are orthonormal vectors, $n = 2d$, and
\[S = (-\mu \alpha \otimes \beta, \text{Id}) \in \mathbb{R}^{d \times 2d}.
\]

As $\text{ker} S = \text{span}\{(\alpha, 0), (\beta, \mu\alpha)\}$ any $S$-objective function $f$ satisfies
\[f(z, w) = f(z + x\alpha + y\beta, w + \mu y\alpha) \text{ for all } x, y \in \mathbb{R}.
\]
Moreover, by Proposition 2 part 2
\[
\nabla_z f \cdot \alpha = 0,
\]
\[
\nabla_z f \cdot \beta + \mu \nabla_w f \cdot \alpha = 0
\]

or, equivalently,
\[\nabla_z f = -\mu (\nabla_w f \cdot \alpha) \beta.
\]

Our results are based on the observation that the representation of objective functions as in Proposition 2 is not unique because $S$ is not fully determined by the null space. A careful choice of the representation can lead to interesting results.

**Definition 3.** Let $S \in \mathbb{R}^{n \times d}$ be a matrix. A function $f$ is rescaled $S$-objective if it admits the representation
\[f(\xi) = \det \eta G(\eta S \xi)
\]
for some density $G$, where $\eta \in \mathbb{R}^{d \times d}_{\text{sym}}$.

In the shear flow setting one obtains the scaling relation
\[p = \eta (w + \mu \alpha \otimes \beta z),
\]
and the corresponding differential relation
\[\nabla_w f = \eta \nabla_p G.
\]

Rescaled solutions for the Boltzmann equation in shear flow settings have been considered in numerous publications, in particular, [10] and [20]. Our main contribution to this topic is the consideration of a renormalization operator $\eta$ which is nonisotropic, i.e., $\eta \neq \lambda \text{Id}$ for all $\lambda \in \mathbb{R}$.
3. The Fokker–Planck case. The Fokker–Planck equation is typically considered as the Kolmogorov forward equation of a Brownian particle in a fluid. It has also been proposed as an approximation of the Boltzmann equation, e.g., in [26, 9]. Furthermore [22, 18] use Fokker–Planck equations to study grazing collisions in the Boltzmann equation and the Kac model. Carlen and Gangbo use a Fokker–Planck equation also as a model problem in [5] for the descent in a Wasserstein metric in kinetic equations; further extensions are given in [6].

Normally the kinetic energy \( \theta \) is a fixed parameter in the Fokker–Planck equation. In our setting we assume that \( \theta \) depends on the density \( f \), as a result the structural properties of the solutions are very similar to the solutions of the Boltzmann equation. In particular, mass, momentum, and energy are conserved, however energy conservation only holds for \( \mu = 0 \). Letting \( \xi = (z, w) \in \mathbb{R}^{2d} \),

\[
\begin{cases}
\partial_t f_t(\xi) = L f_t(\xi), & \xi \in \mathbb{R}^{2d}, t > 0, \\
f_0(\xi) = g_0(S \xi), & \xi \in \mathbb{R}^{2d}, t = 0,
\end{cases}
\]

with \( g_0 \in L^1(\mathbb{R}^d), g_0 \geq 0 \),

\[
Lf(z, w) = -w \cdot \nabla z f(z, w) + \Delta_w f(z, w) + \frac{\rho(z)}{2\theta(z)} \nabla_w \cdot (f(z, w) \left( w - \frac{1}{\rho(z)} v(z) \right)),
\]

and thermodynamic quantities depending on the space variable \( z \),

\[
\begin{align*}
\rho(z) &= \int f(z, w) \, dw \quad \text{(density)}, \\
\theta(z) &= \frac{1}{2} \int |w - \rho^{-1} v(z)|^2 f(z, w) \, dw \quad \text{(kinetic energy)}, \\
v(z) &= \int w f(z, w) \, dw \quad \text{(momentum)}.
\end{align*}
\]

For the solutions of interests, integration over \( z \) will not lead to finite quantities. However the motivation for (6) is that it is similar to the classical Boltzmann equation as it has comparable conservation properties. To see this we define for \( S \)-objective solutions the standard thermodynamic quantities, which can depend on time along a solution \( f_t \), by evaluating at \( z = 0 \)

\[
m_t = \rho_t(0) \text{(mass)}, \quad \bar{v}_t = v_t(0) \text{(momentum)}, \quad \theta_t = \theta_t(0) \text{(energy)}.
\]

The values for other \( z \) are then determined by objectivity.

**Proposition 4.** Let \( d = 2 \) and let \( f \) be a solution of (6) and (1) such that \( \sup_{0 \leq t < T} \theta[f_t] < \infty \). Then there exists \( g_t \) such that \( f_t(\xi) = g_t(S \xi) \). Furthermore, mass \( m \) and \( \bar{v} \) are conserved.

If \( \mu = 0 \), then energy \( \theta \) is also conserved. If \( f^M \) is a Maxwellian, i.e., \( f^M(w) = \exp(h(w)) \) and

\[
h(w) = a + b \cdot w + c |w|^2
\]

for some \( a \in \mathbb{R}, b \in \mathbb{R}^2, c < 0 \), then \( f \) is a stationary solution. Any spatially homogenous \( f \) with \( f(.) (1 + |z|^2) \in L^1(\mathbb{R}^2) \) converges to some \( f^M \) with an exponential rate as \( t \to \infty \).

The existence of \( g_t \) immediately follows from Proposition 2. The rest of the proof mainly involves direct calculations, which we postpone to the appendix.
Equations (6) and (1) do not admit stationary solutions if $\mu \neq 0$. We now aim to characterize the asymptotic behavior of objective solutions for nonzero $\mu$. The main result of this section states that there exists a time-dependent rescaling operator $\eta_t$ such that $G_t$ converges to a Maxwellian as $t \to \infty$.

**Theorem 5.** Let $d = 2$. There exists $\eta_t \in C^1([0, \infty), \mathbb{R}_{\text{sym}}^{2 \times 2})$ such that the rescaled Fokker–Planck equation

\begin{equation}
\begin{aligned}
\partial_t G_t &= \nabla \cdot \left( G_t(p) \left( \theta_t^{-1} \text{Id} - F_t \right) p + \eta_t^2 \nabla G_t \right), \quad t > 0, \ p \in \mathbb{R}^2, \\
F_t &= (\eta_t - \mu \eta_t \alpha \otimes \beta) \eta_t^{-1}, \quad t > 0,
\end{aligned}
\end{equation}

admits a global solution $G_t$ if $G_0 \in L^1 \cap L^\infty$ and $\int_{\mathbb{R}^2} G_0(p)(1 + |p|^2) \, dp < \infty$. The density $f$, which is defined by (3), satisfies (6) and (1).

Furthermore, assume $\int_{\mathbb{R}^2} G_0(p) \, dp = 1$ and $\int_{\mathbb{R}^2} G_0(p) p \, dp = 0$. The density $G_t$ converges to the Maxwellian $G^M(p) = (4\pi)^{-\frac{d}{2}} \exp(-\frac{1}{2}|p|^2)$ for large $t$ in the $L^1$ sense with an algebraic rate, i.e., there exist $\lambda_-, \lambda_+ > 0$ such that

\begin{equation}
\limsup_{t \to \infty} t^{\lambda_-} \|G_t - G^M\|_{L^1(\mathbb{R}^2)} < \infty \quad \text{for all } G_0
\end{equation}

and

\begin{equation}
\liminf_{t \to \infty} t^{\lambda_+} \|G_t - G^M\|_{L^1(\mathbb{R}^2)} > 0
\end{equation}

for $G_0$ in an open dense set of admissible initial data with $\int_{\mathbb{R}^2} G_0(p)(1 + |p|^6) \, dp < \infty$.

Furthermore, for $t \to \infty$ the rescaling operator $\eta_t$ admits the asymptotics:

\begin{equation}
\eta_t = \frac{1}{\mu t^2 + O(t)} \left( \sqrt{3} \alpha \otimes \alpha + 3(\alpha \otimes \beta + \beta \otimes \alpha) + 2\mu t \beta \otimes \beta \right).
\end{equation}

**Remark 6.**

1. The assumption that $d = 2$ is not necessary. The same result can be obtained if $d \geq 2$ at the expense of more complicated notation.
2. If $\mu \neq 0$ the energy is not conserved. As a result (6) is nonlinear, hence, even long-time existence and uniqueness of solutions is not completely trivial.
3. The fact that Maxweilians are global attractors of the dynamics is typically attributed to the observation that the entropy is a Lyapunov functional. We show in the appendix that the functional

\begin{equation}
S[G] = \int_{\mathbb{R}^2} \left( \log G(p) + \frac{1}{2} |p|^2 \right) G(p) \, dp
\end{equation}

decreases for solutions of the Fokker–Planck equation under shear $S$-objectivity. However, this observation is not sufficient for the solutions to converge to the minimum of $S$ as the dissipation operator $\nabla \cdot (\eta_t^2 \nabla \cdot)$ degenerates for $t \to \infty$ as in (10).

4. The algebraic order $\lambda_- > 0$ follows from Proposition 11 below. It is not explicit as we use an abstract result of [33] to obtain it. Similarly, our calculation of the constant for the lower bound $\lambda_+$ is relatively crude. The lower algebraic estimates are based on a detailed understanding of fourth and sixth order moment equations. The analysis provides lower estimates for all such data, which have—after rescaling with $\eta_t$—different fourth moments compared to $G^M$. However, our method does not provide explicit estimates if the fourth moments of the initial distribution and the corresponding Maxwellian coincide.
5. We can also consider general \( \int_{\mathbb{R}^2} G_0(p) \, dp = m > 0 \) and \( \int_{\mathbb{R}^2} G_0(p)p \, dp = mv \in \mathbb{R}^2 \). A translation of the coordinate system can remove the drift, the different mass will need to be introduced in the normalization condition (16) for \( \eta_t \) below. Then \( G \) will converge to \( mG^M(\cdot, +v) \).

The proof will take up the rest of this section. It involves several steps.

1. In the beginning of subsection 3.1 we derive a differential equation for the representative \( g \) of the \( S \)-objective function \( f \) and construct a solution to this equation in Proposition 7.

2. In subsection 3.2 we define the rescaling operator \( \eta_t \) and the shape \( G \). Their asymptotic behavior is obtained from a closed system of moment equations as stated in Proposition 8.

3. The main ingredients of the convergence proof are given in subsection 3.3. We show that the Maxwellian \( G^M \) is an equilibrium and use Proposition 8 to identify the leading terms in (7). After an appropriate rescaling of time the equation has the form of an autonomous degenerate parabolic part plus small nonautonomous perturbations. Hypocoercivity estimates are used for the autonomous degenerate parabolic part in \( H^1 \) relative to the Maxwellian. Additional a priori estimates for the full equation in higher Sobolev norms are provided using calculations inspired by the hypocoercivity framework. The convergence results follow with a Duhamel formula.

4. The equations for fourth and sixth moments are used in subsection 3.4 to obtain the lower estimates.

5. The proof is summarized in subsection 3.5.

### 3.1. Reformulation and regularity.

To minimize the notation we will assume that \( \int G_0(p) \, dp = 1 \) and \( \int G_0(p)p \, dp = 0 \), i.e., \( m = 1 \) and \( \bar{v} = 0 \).

If \( f_t \) is an objective solution, i.e., \( f_t(z, w) = g_t(w + \mu \alpha \otimes \beta z) \), then by (6) and (2) \( g_t \) satisfies

\[
\partial_t g = \mu (\nabla_w g \cdot \alpha) (\beta \cdot w) + \Delta_w g + \theta^{-1} \nabla_w \cdot (g(w) \, w)
\]

\[
= \nabla_w \cdot \left( (g(w) (\theta^{-1} \Id + \mu \alpha \otimes \beta) \, w) + \Delta_w g \right).
\]

A rescaled objective solution \( G_t(p) = \det \eta_t^{-1} g_t(\eta_t^{-1} p) \) with \( p = \eta_t w \) satisfies

\[
\partial_t G_t(p) = \partial_t \left( g(\eta_t^{-1} p) \, \det \eta_t^{-1} \right)
\]

\[
= (\partial_t g - \nabla_w g : \eta_t^{-1} \eta_t^{-1} - \Id) g \, \det \eta_t^{-1}
\]

\[
= (\partial_t g - \nabla_w \cdot (g(\eta_t^{-1} w))) \, \det \eta_t^{-1}
\]

\[
= (\nabla_w \cdot (g(\eta_t^{-1} p) (\theta^{-1} \Id + \mu \alpha \otimes \beta) \, w) + \Delta_w g - \nabla_w \cdot (g(\eta_t^{-1} \eta_t w))) \, \det \eta_t^{-1}.
\]

Now observe that \( \nabla_w = \eta_t \nabla_p \). Continuing the above calculation we obtain

\[
\partial_t G_t = (\nabla_p \cdot \eta_t (G_t(p) (\theta^{-1} \Id + \mu \alpha \otimes \beta) \eta_t^{-1} p) + \nabla_p (g_t(\eta_t \eta_t^{-1} p))) \, \det \eta_t^{-1}
\]

\[
+ \nabla_p \cdot \eta_t^2 \nabla_p G_t = \nabla_p = (G_t(p) (\theta^{-1} \Id - (\eta_t - \mu \eta_t \alpha \otimes \beta) \eta_t^{-1} p) + \eta_t^2 \nabla_p G_t(p))
\]

which is (7).
Next we show that (7) admits unique solutions for arbitrary times. It suffices to consider the case $\eta_t = \text{Id}$ because for a general function $\eta_t \in C^1([0, \infty), \mathbb{R}^{2 \times 2})$ the density $G(p) = \det \eta_t^{-1} g(\eta_t^{-1} p)$ satisfies (7) if $g$ solves (12).

We formulate the underlying regularity result next. The diffusion term $\Delta_w g$ is the generator of the strongly continuous semigroup on $L^2(\mathbb{R}^2)$ via convolution with the classical heat kernel $\Phi_t(\cdot) = \frac{1}{\sqrt{4\pi t}} \exp(-\frac{|\cdot|^2}{4t})$ for $t > 0$. Following, e.g., [25], the kernel $\Phi_t(\cdot)$ also generates equispectral semigroups on weighted $L^p$ spaces like $L^2_2(\mathbb{R}^2)$, i.e., the space of integrable functions that satisfy $\int_{\mathbb{R}^2} g_0^2(w) (1 + |w|^2) \, dw < \infty$.

Due to $\theta$ (12) is nonlinear in $g$, furthermore, the factor $w$ makes the divergence terms unbounded. Hence we need to take care to define a mild solution to (12) to be a solution in $L^2_2(\mathbb{R}^2)$ of the form

$$g_t = [\Phi_t * g_0] + \int_0^t [\Phi_{t-s}(* \cdot \nabla \cdot (g_s(\cdot) \alpha \otimes \beta)) + \theta^{-1} \nabla \cdot (g_s(\cdot) \cdot \theta_s)] \, ds.$$ \hspace{1cm} (13)

**Proposition 7.** Let $g_0 \in L^1(\mathbb{R}^2) \cap L^\infty(\mathbb{R}^2)$. If $\int_{\mathbb{R}^2} g_0(w) (1 + |w|^2) \, dw < \infty$, then (12) admits a unique mild solution for all $t > 0$ such that $\int_{\mathbb{R}^2} g_t(w) (1 + |w|^2) \, dw < \infty$. Furthermore $g_t$ is smooth for $t > 0$.

**Proof.** By Hölder’s inequality, we have $g_0 \in L^2_2(\mathbb{R}^2)$, such that the first term $\Phi_t * g_0$ in (13) is well-defined. We will obtain a mild solution as in (13) via an approximation scheme using nonautonomous bounded perturbations. Letting

$$\chi_n(w) = \begin{cases} w & \text{if } |w| < n, \\ n |w|/|w| & \text{otherwise,} \end{cases}$$

and $\theta_0 = 1$, we define recursively for $n \in \mathbb{N}$ as a nonautonomous Miyarawa perturbation (see, e.g., [32]), the following mild solution

$$g^n_t = [\Phi_t * g_0] + \int_0^t [\Phi_{t-s}(* \cdot \nabla \cdot (g^n_s(\cdot) \alpha \otimes \beta \chi_n(\cdot)) + \theta^{-1}_{n-1} \nabla \cdot (g^n_s(\cdot) \cdot \theta_s)] \, ds.$$ \hspace{1cm} (14)

By the properties of the convolution, we see that $g^n$ is smooth for $t > 0$ and gives a classical solution as the second convolution in (14) is well-defined. For a fixed time $T > 0$ standard a priori estimates give uniform bounds in $L^\infty((0, T), L^2_2(\mathbb{R}^2))$ and $L^2((0, T), H^1_2(\mathbb{R}^2))^*$.

Differentiating $\theta_n$ with respect to $t$ gives

$$\frac{d\theta_n}{dt} = \frac{1}{2} \int_{\mathbb{R}^2} |w|^2 (\cdot \nabla \cdot (g^n_t(\cdot) \alpha \otimes \beta w) + \Delta g^n_t + \theta^{-1}_{n-1} \nabla \cdot (g^n_t w)) \, dw$$

$$= \frac{1}{2} \int_{\mathbb{R}^2} |w|^2 \Delta g^n_t \, dw - \mu \int_{\mathbb{R}^2} (w \cdot \alpha)(\beta \cdot w) g^n_t \, dw - \theta^{-1}_{n-1} \int_{\mathbb{R}^2} |w|^2 g^n_t \, dw$$

$$\leq C + (\frac{2}{\theta_{n-1}}) \theta_n.$$  

Thus with a Gronwall estimate, $\theta_n$ and $\theta_n$ remain bounded. Similarly,

$$\frac{d}{dt} (\theta_n)^{-1} = -\frac{d\theta_n}{dt} \frac{1}{\theta_n^2} = -\frac{1}{\theta_n^2} \left( 2 - \mu \int_{\mathbb{R}^2} (w \cdot \alpha)(\beta \cdot w) g^n_t \, dw - \theta^{-1}_{n-1} 2 \theta_n \right)$$

$$\leq -\frac{2}{\theta_n^2} + \left( \frac{2}{\theta_{n-1}} \right) (\theta_n)^{-1}$$

$$= -\frac{2}{\theta_n} \left( \frac{1}{\theta_n} - \frac{1}{\theta_{n-1}} \right) + \frac{2}{\theta_{n-1}} \theta_n^{-1},$$  

(15)
which also shows that $\theta_n^{-1}$ and $\frac{d}{dt}(\theta_n)^{-1}$ remain bounded on $(0, T)$ as for fixed $n$ the penultimate line gives that $\theta_n^{-1}$ cannot grow beyond $|\mu|/2 + \theta_n^{-1}$. We obtain the bound $\theta_n^{-1}(t) \leq \frac{1}{\delta_0} \exp(|\mu|t)$ for $t \in (0, T)$ uniformly in $n$ by induction: For $n = 0$, $\theta_0^{-1}$ is constant and trivially $\theta_0^{-1}(t) \leq \frac{1}{\delta_0} \exp(|\mu|t)$ for $t \in (0, T)$. Now suppose the estimate holds up to $n - 1$. If $\theta_n^{-1}(t) \geq \frac{1}{\delta_0} \exp(|\mu|t)$, then $\theta_n^{-1} \geq \theta_{n-1}^{-1}$, such that the nonlinear term is negative in the last line of (15), such that $\frac{d}{dt}(\theta_n)^{-1} \leq |\mu|(\theta_n)^{-1}$.

This implies the Gronwall estimate $\theta_n^{-1}(t) \leq \frac{1}{\delta_0} \exp(|\mu|t)$ for $t \in (0, T)$.

All bounds combined give a subsequence, which we do not relabel, such that $g_n \to g$ weakly in $L^2((0, T), H^1_2(\mathbb{R}^2))$, $g_n \to g$ weak star in $L^\infty((0, T), L^2_2(\mathbb{R}^2))$, $\theta_n \to \theta$ strongly in $C^0(0, T)$ by Arzela–Ascoli, such that $1/\theta$ is also bounded. Then the nonlinear term will converge weakly to $\theta^{-1} \nabla \cdot (g(\cdot, \cdot))$ in $L^2((0, T), L^2_2(\mathbb{R}^2))$. Next we show that $g$ satisfies (13). We first observe that when the integral in the variation-of-constants formula is restricted to

$$\int_0^{t-\delta} \Phi_{t-s}(\cdot) \ast (\mu \nabla \cdot (g^n(\cdot) \alpha \otimes \beta) + \theta^{-1}_{n-1} \nabla \cdot (g^n(\cdot)) \, ds$$

for any $\delta > 0$ then weak convergence in $L^2((0, T), L^2(\mathbb{R}^2))$ is enough to show its convergence to

$$\int_0^{t-\delta} \Phi_{t-s}(\cdot) \ast (\mu \nabla \cdot (g(\cdot) \alpha \otimes \beta) + \theta^{-1} \nabla \cdot (g(\cdot)) \, ds,$$

as $\Phi_{t-s}(\cdot) \in L^2((0, t-\delta) \times \mathbb{R}^2)$. The remainder term is $O(\sqrt{\delta})$, which ensures convergence, such that the weak limit of the subsequence $g^n$ satisfies (13), which implies differentiability for $t > 0$. To show uniqueness consider two solutions $g, h$ with $\theta|g|^{-1}$ and $\theta|h|^{-1}$ bounded for the same initial data. We obtain the inequalities

$$\frac{d}{dt} \|g(t) - h(t)\|_{L^2_2} \leq C|\theta(g(t))^{-1}|\|g(t) - h(t)\|_{L^2_2} + C|\theta(g(t))^{-1} - \theta(h(t))^{-1}|,$$

$$\frac{d}{dt} (\theta|g(t)|^{-1} - \theta|h(t)|^{-1}) \leq C\|g(t) - h(t)\|_{L^2_2} + C|\theta(g(t))^{-1} - \theta(h(t))^{-1}|,$$

which show $g = h$ using the Gronwall inequality. This also yields that there is a unique limit for the weakly convergent subsequence.

\[\square\]

### 3.2. Moment equations.

The next step is to study the evolution equations of the moments of $g$ and $G$.

A careful analysis of the moments of $g$ and $G$ will deliver the following.

1. The rescaling operator $\eta_t$ by requiring that

$$\frac{1}{2} \int G(p) p \otimes p \, dp = \text{Id}$$

holds for all $t \geq 0$.

2. Tightness of $p^2 G(p)$.

An easy calculation shows that (16) holds if

$$\eta_t = T^{-\frac{1}{2}},$$

where

$$T = \frac{1}{2} \int g_t(w) w \otimes w \, dw$$
is the Cauchy stress tensor for \( g \). Indeed,
\[
\frac{1}{2} \int G(p) p \otimes p \, dp = \frac{\det \eta_t}{2} \int (\eta_t \omega) \otimes (\eta_t \omega) G(\eta_t \omega) \, d\omega = \eta_t T \eta_t^* \]
as required.

Finally we characterize the long-time behavior of \( G_t \) if \( \eta_t = T^{-\frac{1}{2}} \). The results are summarized in the next proposition.

**Proposition 8.** Letting \( G \) be a solution of (7) with initial data as in Theorem 5 and \( \eta_t = T^{-\frac{1}{2}} \), then the following asymptotics hold for \( t \to \infty \):

\[
\begin{align*}
\frac{\eta_t}{t^{\frac{3}{2}} + O(t)} &= \frac{1}{\mu (t^{\frac{1}{2}} + O(t))} (\sqrt{3} t\alpha \otimes \alpha + 3(\alpha \otimes \beta + \beta \otimes \alpha) + 2\mu t \beta \otimes \beta), \\
T^{-1} &= \frac{2}{t + O(t^2)} \left( \frac{6}{(\mu t)^2} \alpha \otimes \alpha + \frac{3}{\mu t} (\alpha \otimes \beta + \beta \otimes \alpha) + 2\beta \otimes \beta \right), \\
\theta^{-1} &= O(t^{-3}), \\
F &= (\eta_t - \eta_t \alpha \otimes \beta) \eta_t^{-1} \\
&= -\frac{1}{2t + O(t)} \left[ O(1/t^2) \alpha \otimes \alpha + \sqrt{3}(\alpha \otimes \beta - \beta \otimes \alpha) + 4\beta \otimes \beta \right].
\end{align*}
\]

Furthermore there exist \( c, \lambda, \lambda' > 0 \) such that for all permissible \( G_0 \in L^1_0(\mathbb{R}^2) \) we have that
\[
\int_{\mathbb{R}^2} |(G_t - G^M)(p)| \, |p|^6 \, dp = O(1 + t^{\lambda'})
\]
and for an open dense set of initial data \( G_0 \in L^1_0(\mathbb{R}^2) \) there exists \( c > 0 \) such that for sufficiently large \( t \)
\[
\left| \int_{\mathbb{R}^2} (G_t - G^M)(p) \, |p|^4 \, dp \right| > ct^{-\lambda}.
\]

**Proof.** We will first establish formulas (18), (19), (20), and (21) by carefully analyzing the second moments. Formulas (22) and (23) follow from cruder estimates of higher moments.

**Second moments.** The stress tensor \( T = \frac{1}{2} \int g_t(\omega) \omega \otimes \omega \, d\omega \) satisfies a simple ordinary differential equation. Multiplying (12) with \( \frac{1}{2} \omega \otimes \omega \) and integrating by parts yields
\[
\frac{dT}{dt} = \text{Id} - \mu (\alpha \otimes \beta T + T \beta \otimes \alpha) - \frac{2}{\mu T} T.
\]
To characterize the asymptotic behavior of \( T \) as \( t \to \infty \) we define the rescaled moments \( a, b, c \) by the requirement
\[
T = t^3 a \alpha \otimes \alpha + t^2 b (\beta \otimes \alpha + \alpha \otimes \beta) + tc \beta \otimes \beta.
\]
Then (24) reads
\[
3 t^2 a \alpha \otimes \alpha + 2t b (\beta \otimes \alpha + \alpha \otimes \beta) + c \beta \otimes \beta + t^3 \frac{da}{dt} \alpha \otimes \alpha + t^2 \frac{db}{dt} (\beta \otimes \alpha + \alpha \otimes \beta) + t \frac{dc}{dt} \beta \otimes \beta
\]
\[
= \alpha \otimes \alpha + \beta \otimes \beta - \mu (2t^2 b \alpha \otimes \alpha + tc (\alpha \otimes \beta + \beta \otimes \alpha))
\]
\[
- \frac{2}{a t^3 + ct} \left( t^3 a \alpha \otimes \alpha + t^2 b (\beta \otimes \alpha + \alpha \otimes \beta) + tc \beta \otimes \beta \right).
\]
The equations for the individual components read
\begin{align*}
\bfalpha \otimes \bfalpha : 0 &= 3t^2a + t^3 \frac{da}{dt} - 1 + 2t^2 \mu b + \frac{2at^3}{at^3 + ct}, \\
\bfalpha \otimes \bfbeta : 0 &= 2tb + t^2 \frac{db}{dt} + \mu tc + \frac{2t^2b}{at^3 + ct}, \\
\bfbeta \otimes \bfbeta : 0 &= c + t \frac{dc}{dt} - 1 + \frac{2ct}{at^3 + ct}.
\end{align*}

After rescaling time as well so that \(t = \exp(s)\) and \(\frac{dt}{ds} = \exp(-s) \frac{ds}{dt}\), (24) takes the form
\[
\left( \frac{d}{ds} - M \right) \begin{pmatrix} a \\ b \\ c \end{pmatrix} = \begin{pmatrix} \exp(-2s) \\ 0 \\ 1 \end{pmatrix} + \frac{2}{\exp(2s) a + c} \begin{pmatrix} a \\ b \\ c \end{pmatrix}
\]
with
\[
M = \begin{pmatrix} -3 & -2\mu & 0 \\ 0 & -2 & -\mu \\ 0 & 0 & -1 \end{pmatrix}.
\]

As the spectrum of \(M\) is given by \(\lambda_1 = -1, \lambda_2 = -2, \lambda_3 = -3\) with corresponding eigenvectors \(v_1 = (1, 0, 0), v_2 = (-2\mu, 1, 0), v_3 = (\mu^2, -\mu, 1)\), a simple application of the variation of constants formula delivers the asymptotic result
\[
\begin{pmatrix} a \\ b \\ c \end{pmatrix} = M^{-1} \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} + O(\exp(-s)) = \begin{pmatrix} \frac{1}{2} \mu^2 \\ -\frac{1}{2} \mu \\ 1 \end{pmatrix} + O(t^{-1}).
\]

This implies that the stress tensor admits the asymptotic result
\[
T = (t + O(1)) T_\infty, \quad t \to \infty,
\]
where
\[
T_\infty = \frac{1}{3} (t\mu)^2 \bfalpha \otimes \bfalpha - \frac{1}{2} t\mu (\bfalpha \otimes \bfbeta + \bfbeta \otimes \bfalpha) + \bfbeta \otimes \bfbeta.
\]

We can bootstrap this step by plugging (25) into (24). This shows that the function \(t \to T - t T_\infty\) is differentiable and satisfies
\[
\frac{d}{dt} (T - t T_\infty) = O(t^{-1}), \quad t \to \infty.
\]

These asymptotics give results for \(\eta, \eta^{-1}, \dot{\eta},\) and \(T^{-1}\). Then as \(\theta^{-1} = (\text{tr} \, T)^{-1}\), this implies (19) and one also has the asymptotic result for \(F\).

**Fourth and sixth moments.** We are now assuming that the related initial data for \(G_0\) satisfy \(g_0 \in L^6_0(\mathbb{R}^2)\) and using (13) we can see that higher moments up to order 6 are well-defined for finite times; these moments satisfy similar ordinary differential equations. These will later allow us to choose suitable initial data for lower estimates on the rate of decay. Letting
\[
h_{ij}(t) = \int_{\mathbb{R}^2} (G_t(p) - G^M(p)) p_i^1 p_j^2 \, dp,
\]
(27)
we obtain ordinary differential equations, which only depend on modes of the same or lower order. The moment of order 0 and 1 (mass and momentum) are preserved by Proposition 4 for the evolution of $g$, in the same way this also follows for $G$, where the momentum is assumed to be 0, i.e.,

$$h_{00}(t) \equiv 0 = h_{10}(t) \equiv h_{01}(t) \equiv 0.$$  

(28)

The rescaling $\eta_t$ is defined such that (16) holds, i.e.,

$$h_{20}(t) \equiv h_{02}(t) \equiv h_{11}(t) \equiv 0.$$  

(29)

For the higher moments $h_{ij}$ with $i + j > 2$ we obtain using integration by parts

$$\frac{d}{dt} h_{ij}(t) = -\left(\frac{i + j}{\theta} - iF_{11} - jF_{22}\right) h_{ij} + jF_{12} h_{i+1,j-1} + iF_{21} h_{i-1,j+1}$$

$$+ (i - 1)(\eta_t^2)_{11} h_{i-2,j} + ij((\eta_t^2)_{12} + (\eta_t^2)_{21}) h_{i-1,j-1} + j(j - 1)(\eta_t^2)_{22} h_{i,j-2},$$

where we assumed without loss of generality that $\alpha = (1, 0)$ and $\beta = (0, 1)$. Using the information on the coefficients in Proposition 8 we write the moment equations in matrix notation:

$$\frac{d}{dt} h_{ij}(t) = t^{-1} \sum_{k,l} (N_{ijkl} + O(t^{-1})) h_{kl}(t) \quad t \gg 1,$$

(30)

where the operator $N$ is defined by

$$N_{ijkl} = \begin{cases} 
-2j & \text{if } (i, j) = (k, l), \\
-\sqrt{3}j & \text{if } (i, j) = (k + 1, l - 1), \\
\sqrt{3}j & \text{if } (i, j) = (k - 1, l + 1), \\
2j(j + 1) & \text{if } (i, j) = (k, l + 2), \\
0 & \text{else},
\end{cases}$$

with the convention that $h_{ij} = 0$ if $i + j \leq 2$. The moments of odd order can all be chosen to be 0, which is preserved by (30). Now rescaling time $t = \exp(s)$ and $\frac{d}{dt} = \exp(-s) \frac{d}{ds}$, (31) becomes

$$\left(\frac{d}{ds} - N + O(e^{-s})\right) h = 0, \quad s \gg 1,$$

(32)

where the operator $N$ is a lower triangular form. Consider now the truncated operator

$$(N_{ijkl})_{i+j=k+l=4} = \begin{pmatrix} 
0 & 2\sqrt{3} & 0 & 0 & 0 \\
-\sqrt{3}/2 & -2 & 3\sqrt{3}/2 & 0 & 0 \\
0 & -\sqrt{3} & -4 & \sqrt{3} & 0 \\
0 & 0 & -3\sqrt{3}/2 & -6 & \sqrt{3}/2 \\
0 & 0 & 0 & -2\sqrt{3} & -8
\end{pmatrix}.$$

It has only eigenvalues with negative real parts by the Routh–Hurwitz stability criterion. Using the variation of constants formula we obtain
\[(h_{ij})_{i+j=4} = O(e^{-\lambda s})\]
for some \(\lambda > 0\). Letting \(u(s) = \exp(-Ns)h(s)\) changes (32) into
\[
\left( \frac{d}{ds} + O(e^{-s}) \right) u = 0.
\]
As the \(O(e^{-s})\) term is integrable, bounded initial data \(u(0)\) will remain bounded in the norm from above and below for all times \(s > 0\). Hence there is some \(\lambda \geq \tilde{\lambda}\),
\[
\liminf_{s \to \infty} \exp(\lambda s) \left| (h_{ij})_{i+j=4} \right| > 0
\]
for all nonzero initial data in (30). All initial data \(G_0 \in L^1(R^2)\) with a different tensor of fourth order moments—after the coordinate change \(\eta_0\)—compared to \(G^M\) will then have nonzero initial data in (30); this set is open and dense in the set of possible initial data in \(L^1(R^2)\). Transferring this back to time \(t\) gives the algebraic estimate (23) for some \(c > 0\) which depends on the initial value and all \(t\) large enough.

A less detailed calculation for the vector \(h\) of sixth moments then gives
\[
\frac{d}{dt} h = N(6)(t)h + O \left( \frac{1}{1 + t^\alpha} \right),
\]
where \(\alpha > 1\) and \(\|N(6)(t)\| = O(\frac{1}{t^\alpha})\) as \(t \to \infty\). After transforming to time \(s\) as above we obtain a constant matrix plus some exponentially small error terms; this is enough using Gronwall’s inequality to conclude that \(h\) grows at most exponentially in \(s\) and after transforming back to time \(t\) that \(h\) grows at most algebraically with rate \(t^{\lambda'}\) such that—without loss of generality—\(\lambda' \geq 0\), this then yields (22).

3.3. Asymptotics of shape equation and hypocoercivity. With the asymptotic information on the coefficients we can study the shape equation (7).

**Lemma 9.** The Maxwellian \(G^M(p) = \frac{1}{4\pi} \exp(-\frac{1}{4}|p|^2)\) is a stationary solution to (7).

*Proof.* Substituting \(G^M\) into (7) and observing that \(\nabla G^M(p) = -\frac{1}{2} G^M(p) p\) one finds that \(G^M\) is stationary if and only if
\[
\text{tr} \left( \theta^{-1} \text{Id} - F - \frac{1}{2} \eta_2 \right) G^M(p) - \frac{1}{2} p \cdot \left( \theta^{-1} \text{Id} - \frac{1}{2} (F + F^*) - \frac{1}{2} \eta_2 \right) p G^M(p) = 0.
\]
Next, recall that by (16) the covariance matrix of \(G\) is constant, i.e., \(\frac{d}{dt} \int p \otimes p G(p) \, dp = 0\). Hence, after multiplication of (7) with \(\frac{1}{2} p \otimes p\) and integration by parts one obtains that
\[
F + F^* + \eta_2^2 = 2 \theta^{-1} \text{Id}.
\]
Clearly both terms on the left-hand side (35) vanish thanks to (36).
coercivity constant of the dissipation operator $\nabla \cdot \eta_t^2 \nabla$ diverges as $t \to \infty$. To show the convergence we use a nonautonomous perturbation to the theory of hypocoercivity [33] combined with a priori estimates for the full equations.

It is advantageous to rescale time by defining

$$\tilde{G}_s = G_{\exp(s)}.$$  

As $t = \exp s$ and $\frac{dt}{ds} = t$ the density $\tilde{G}$ satisfies the rescaled equation

$$t \partial_t \tilde{G} = \partial_s \tilde{G} = t \nabla \cdot \left( \tilde{G}(p) \left( \theta^{-1} \text{Id} - F \right) p + \eta_t^2 \nabla \tilde{G} \right).$$

The coefficients in (37) are controlled by Proposition 8. Next we rewrite (37) as a density with respect to $G^M$, then we obtain for $G_t = u_t G^M$,

$$\partial_s u = \nabla_p \cdot \left( T^{-1} \nabla u \right) + \nabla u \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) p.$$

We split the last equation into an autonomous main part using (19), (20), and (21) that provide bounds on some decaying perturbation:

$$\partial_s u = 4 \text{tr} (\beta \otimes \beta \nabla^2 u) + \nabla u \cdot \left( \frac{\sqrt{3}}{2} (\alpha \otimes \beta - \beta \otimes \alpha) - 2 \beta \otimes \beta \right) p + \exp(-s) \left\{ \nabla u \cdot \left( C_1 \alpha \otimes \beta + C_2 \beta \otimes \alpha + C_3 \beta \otimes \beta + C_4 \alpha \otimes \alpha \right) p + \left( C_5 (\alpha \otimes \beta + \beta \otimes \alpha) + C_6 \beta \otimes \beta + C_7 \exp(-s) (\alpha \otimes \alpha) \nabla^2 u \right) \right\}$$

for some appropriate uniformly bounded nonautonomous coefficients $C_i$ for $i = 0, \ldots, 7$.

For the long-term convergence of solutions of the shape equation we use Villani's concept of hypocoercivity [33]. Consider a separable Hilbert space $\mathcal{H}$ with inner product $\langle \cdot, \cdot \rangle$, which will be $L^2(\mathbb{R}^2, dG^M)$ in our case. Let $A = (A_1, \ldots, A_m)$ be an unbounded operator for some $m \in \mathbb{N}$ with domain $D(A)$ and let $B$ be an unbounded antisymmetric operator with domain $D(B)$. The theory reduces the convergence to equilibrium of the nonsymmetric operator $L = A^* A + B$, which is not coercive in our case, to the study of the symmetric operator $A^* A + C^* C$ using the commutator $C = [A, B]$. Under appropriate conditions this operator is coercive, which then implies convergence in the abstract Sobolev space $\mathcal{H}^1$ with norm $\|h\|_{\mathcal{H}^1}^2 = \langle h, h \rangle + \langle Ah, Ah \rangle + \langle Ch, Ch \rangle$; in our case this will coincide with $H^1(\mathbb{R}^2, dG^M)$. The simplest form of the theory is enough for our example and it is stated next.

**Theorem 10** (see [33, Theorem 18]). With the above notation, consider a linear operator $L = A^* A + B$ with $B$ antisymmetric, and define the commutator $C := [A, B]$. Assume the existence of constants $\alpha, \beta$ such that

1. $A$ and $A^*$ commute with $C$; $A_i$ commutes with each $A_j$;
2. $[A, A^*]$ is $\alpha$-bounded relative to $I$ and $A$;
3. $[B, C]$ is $\beta$-bounded relative to $A, A^2, C$, and $AC$.

Then there is a scalar product $\langle \langle \cdot, \cdot \rangle \rangle$ on $H^1(\mathbb{R}^2, dG^M)/K$, which defines a norm equivalent to the $H^1$ norm, such that

$$\forall h \in H^1/K, \quad \langle \langle h, L h \rangle \rangle \geq K (\|Ah\|^2 + \|Ch\|^2)$$

for some constant $K > 0$ depending on $\alpha$ and $\beta$. If, in addition,

$$A^* A + C^* C$$

is $\kappa$-coercive
for some $\kappa > 0$, then there exists a constant $\lambda > 0$, such that
\[ \forall h \in H^1/K, \quad \langle \langle h, Lh \rangle \rangle \geq \lambda \langle h, h \rangle. \]
In particular, $L$ is hypocoercive in $H^1/K$, there is a $c < 0$,
\[ \| \exp(-tL) \|_{H^1/K \to H^1/K} \leq c \exp(-\lambda t), \]
where both $\lambda$ and $c$ only depend on upper bounds for $\alpha$ and $\beta$ and lower bounds on $\kappa$.

The last theorem is used to show that the leading order of (39), i.e., its autonomous part, is hypocoercive. Then we use a similar splitting $L_s = A_s^* A_s + B_s$ for the full equation to obtain a priori estimates. Both ingredients will then be combined via a Duhamel formula to provide the convergence result for (39).

**Proposition 11.** The autonomous part of (39) given by

\[ \partial_s u = 4 \text{tr} (\beta \otimes \beta \nabla^2 u) + \nabla u \cdot \left( \frac{\sqrt{3}}{2} (\alpha \otimes \beta - \beta \otimes \alpha) - 2 \beta \otimes \beta \right) p \]
defines a contraction in time in $H^1(\mathbb{R}^2, dG^M)/K$, where $K = \text{span}\{1\}$, i.e. there exist $c, \lambda > 0$ such that

\[ \| \exp(-sL) \|_{H^1/K \to H^1/K} \leq c \exp(-\lambda s). \]

**Proof.** Consider $L^2(\mathbb{R}^2, dG^M)$ with inner product $\langle u, v \rangle = \int_{\mathbb{R}^2} u(p) v(p) G^M(p) \, dp$. We are now writing (41) in Villani’s notation

\[ \partial_s u + Lu = 0 \text{ with } L = A^* A + B \text{ and } B \text{ antisymmetric in } L^2(\mathbb{R}^2, dG^M). \]

Choosing coordinates and identifying the canonical basis vectors $e_1, e_2$ with $\alpha$ and $\beta$ respectively, we let

\[ (Au)(p_1, p_2) = 2 \partial_2 u(p_1, p_2), \quad (Bu)(p_1, p_2) = -\frac{\sqrt{3}}{2} (p_2 \partial_1 u(p_1, p_2) - p_1 \partial_2 u(p_1, p_2)). \]

Then we obtain the adjoint $A^*$ of $A$ in $L^2(\mathbb{R}^2, dG^M)$ by integration by parts in the inner product

\[ (A^* u)(p_1, p_2) = -2 \partial_2 u(p_1, p_2) + p_2 u(p_1, p_2), \]

while $B$ is antisymmetric, such that (43) is a reformulation of (41). We now check the assumptions of Theorem 10. We observe $C := [A, B] = -\sqrt{3} \partial_1$ and then (i) $A$ and $A^*$ commute with $C$. Furthermore (ii) holds as $[A, A^*] = 2I$. The commutator $[B, C] = -\frac{3}{2} \partial_2$ is relatively bounded by $A$, hence (iii) holds. The general results then imply $K = \text{Ker} L = \text{Ker} A \cap \text{Ker} B$ consists of constants only. In addition $A^* A + C^* C = -4 \partial_2^2 + 2 p_2 \partial_2 - 3 p_2^2$ is coercive on $L^2(\mathbb{R}^2, dG^M)$ using a Poincaré inequality as in [33, Thm A.1]. Then Theorem 10 implies there exist positive constants $\lambda$ and $c$ such that (42) holds, completing the proof.

To obtain a priori estimates, (39) is rewritten in the form of the last proposition with time-dependent operators $A_s$ and $B_s$:

\[ \partial_s u = -L_s u = -A_s^* A_s u - B_s u, \]
where
\begin{align}
A_s &= \eta_s \nabla, \\
A_s^* &= -\nabla \cdot (\eta_s) + \frac{1}{2} p \eta_s, \\
B_s u &= -\nabla u \cdot \left( \theta^{-1} \text{Id} - F - \frac{1}{2} T^{-1} \right) p.
\end{align}

Then by (35)
\[ B_s^* u = \nabla u \cdot \left( \theta^{-1} \text{Id} - F - \frac{1}{2} T^{-1} \right) p + \frac{1}{5} p \left( \theta^{-1} \text{Id} - F - \frac{1}{2} T^{-1} \right) p = -B_s u, \]
such that \( B_s \) is antisymmetric.

**Lemma 12.** Let \( u_s \) be the solution (39) obtained from rescaling the solution in Proposition 7, then there is \( K_s > 0 \) such the a priori estimates hold for all \( s \geq 1 \):
\[ \| u_s \|_{H^1(\mathbb{R}^2, G_M)} + \| \nabla u_s \|_{H^1(\mathbb{R}^2, G_M)} + \| \nabla^2 u_s \|_{H^1(\mathbb{R}^2, G_M)} \leq K_s. \]

**Proof.** Using the form in (46) we estimate the time derivative of the \( L^2 \) norm
\[ \partial_t (u_s, u_s) = -2 \langle L_s u_s, u_s \rangle = 2 \langle -A_s^* A_s u_s + B_s u_s, u_s \rangle \\
= 2 \langle A_s u_s, A_s u_s \rangle \leq 0. \]

The derivatives of \( \partial_1 u \) and \( \partial_2 u \) with respect to \( p_1 \) and \( p_2 \) satisfy equations similar to (46):
\begin{align}
\partial_s \partial_1 u_s &= -A_s^* A_s \partial_1 u_s - B_s \partial_1 u_s + \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_1, \\
\partial_s \partial_2 u_s &= -A_s^* A_s \partial_2 u_s - B_s \partial_2 u_s + \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_2.
\end{align}

This yields, with the antisymmetry of \( B_s \),
\begin{align*}
\partial_s \langle \nabla u_s, \nabla u_s \rangle &= \partial_s \left( \langle \partial_1 u_s, \partial_1 u_s \rangle + \langle \partial_2 u_s, \partial_2 u_s \rangle \right) \\
&= 2 \langle A_s \partial_1 u_s, A_s \partial_2 u_s \rangle + 2 \langle \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_1, \partial_1 u_s \rangle \\
&+ 2 \langle A_s \partial_2 u_s, A_s \partial_2 u_s \rangle + 2 \langle \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_2, \partial_2 u_s \rangle \\
&\leq 2 \langle \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_1, \partial_1 u_s \rangle \\
&+ \langle \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_2, \partial_2 u_s \rangle \\
&= 2 \langle \nabla u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right), \nabla u_s \rangle \\
&\leq C \exp(-s) \langle \nabla u_s, \nabla u_s \rangle,
\end{align*}
where autonomous terms in (39) either cancel or have a sign; the form of the nonautonomous first-order terms yields the remainder. Then the Gronwall inequality shows that
\[ \langle \nabla u_s, \nabla u_s \rangle \leq \exp \left( C \int_1^s \exp(-\sigma) \, d\sigma \right) \langle \nabla u_1, \nabla u_1 \rangle = \exp \left( C/e \right) \langle \nabla u_1, \nabla u_1 \rangle \]
remains bounded for all times \( s > 1 \). A similar argument also holds for higher derivatives. We derive differential equations for higher derivatives:
\begin{align}
\partial_s \partial_1^2 u_s &= -A_s^* A_s \partial_1^2 u_s - B_s \partial_1^2 u_s + 2 \nabla \partial_1 u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_1, \\
\partial_s \partial_2 \partial_1 u_s &= -A_s^* A_s \partial_2 \partial_1 u_s - B_s \partial_2 \partial_1 u_s + \nabla \partial_1 u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_2 \\
&+ \nabla \partial_2 u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_1, \\
\partial_s \partial_2^2 u_s &= -A_s^* A_s \partial_2^2 u_s - B_s \partial_2^2 u_s + 2 \nabla \partial_2 u_s \cdot \left( \theta^{-1} \text{Id} - F - T^{-1} \right) e_2.
\end{align}
These equations yield, due to the properties of $A_s$ and $B_s$,
\[
\frac{1}{2} \partial_s \left( \langle \partial^2_s u_s, \partial^1_s u_s \rangle + 2 \langle \partial_2 \partial_1 u_s, \partial_1 \partial_1 u_s \rangle + \langle \partial^2_2 u_s, \partial^2_2 u_s \rangle \right)
\leq 2 \langle \nabla \partial_1 u_s \cdot (\theta^{-1} \text{Id} - F - T^{-1}), \nabla \partial_1 u_s \rangle + 2 \langle \nabla \partial_2 u_s \cdot (\theta^{-1} \text{Id} - F - T^{-1}), \nabla \partial_2 u_s \rangle
\leq C \exp(-s) \left( \langle \nabla \partial_1 u_s, \nabla \partial_1 u_s \rangle + \langle \nabla \partial_2 u_s, \nabla \partial_2 u_s \rangle \right),
\]
where the autonomous terms in (39) again either cancel or have a sign; the form of the nonautonomous first-order terms yields the remainder. Using the Gronwall inequality yields a bound on the second derivatives after an initial regularization, e.g., for $s > 1$.

Deriving similar equations for third derivatives and estimating
\[
\partial_s \left( \langle \partial^3_1 u_s, \partial^3_1 u_s \rangle + 3 \langle \partial^2_2 \partial_1 u_s, \partial^2_2 \partial_1 u_s \rangle + 3 \langle \partial_2 \partial^2_1 u_s, \partial_2 \partial^2_1 u_s \rangle + \langle \partial^3_2 u_s, \partial^3_2 u_s \rangle \right)
\]
yields the final required estimate.

It remains to establish the convergence of $G_t$ to $G^M$ in $L^1$. It suffices to show that $u_s \to 0$ in $L^2(\mathbb{R}^2, dG^M)/\text{span}(1)$. Indeed, we show convergence of $u$ in the stronger $H^1(\mathbb{R}^2, dG^M)$ norm. Using the Duhamel principle for the equation,
\[
\partial_s u_s = -Lu_s - (L_s - L)u_s
\]
with $L$ as in (43) and $L_s$ as in (46). We are starting from the positive time 1 for $s > 1$ to guarantee uniform bounds for higher derivatives as in Lemma 12:
\[
u_s = \exp(-L_{s-1}) u_1 - \int_1^s \exp(-L_{s-\sigma})(L_{\sigma} - L) u_{\sigma} \, d\sigma
\]
Then using the error estimates of $L_{\sigma} - L$ in (39), Lemma 12 together with the contraction property of $L$ in $H^1(\mathbb{R}^2, dG^M)$ as in Proposition 11 yields for $s > 1$
\[
\|u_s\|_{H^1(\mathbb{R}^2, dG^M)}
\leq \exp(-\lambda(s - 1)) \|u_1\|_{H^1(\mathbb{R}^2, dG^M)}
+ \int_1^s \| \exp(-L_{s-\sigma})\|_{H^1 \to H^1}(L_{\sigma} - L) u_{\sigma} \|_{H^1(\mathbb{R}^2, dG^M)} \, d\sigma
\leq \exp(-\lambda(s - 1)) \|u_1\|_{H^1(\mathbb{R}^2, dG^M)} + \int_1^s \exp(-\lambda(s - \sigma)) C \exp(-\sigma) K^* \, d\sigma
\leq Cs \exp(-\min\{\lambda, 1\} s)
\]
for some bounded $C$ only depending on the $L^1 \cap L^\infty$ norms of the initial data due to initial regularization. Undoing the change of time from $t$ to $s$ we also see the rate of convergence is bounded by any algebraic order greater than $\min\{\lambda, 1\}$.

3.4. Lower estimates using higher moments. The estimates on the higher moments in the last proposition yield the lower estimates (9) in the following way for almost all initial data.

Letting $B_R$ be the ball of radius $R$ in $\mathbb{R}^2$, we first note that for all $R > 0$ using (22)
\[
\left| \int_{\mathbb{R}^2} (G(p) - G^M(p)) |p|^4 \, dp \right|
\leq R^4 \int_{B_R} |G(p) - G^M(p)| \, dp + R^{-2} \int_{\mathbb{R}^2 \setminus B_R} |G(p) - G^M(p)| |p|^6 \, dp
\leq R^4 \int_{\mathbb{R}^2} |G(p) - G^M(p)| \, dp + \frac{C}{R^2} t'.
\]
Note that $R$ may depend on $t$ in the above estimate.
Then (23) implies with the choice \( R = R(t) = t^{\lambda + \lambda'/2} \) that

\[
\left| \int_{\mathbb{R}^2} (G_t(p) - G^M(p))|p|^4 \, dp \right| \geq 2 \frac{C}{R^2} t^{\lambda}
\]

for \( t \) large enough as the right-hand side is \( O(t^{-2\lambda}) \). Then we obtain that

\[
\liminf_{t \to \infty} t^{5\lambda + 2\lambda'} \| G_t - G^M \|_{L^1(\mathbb{R}^2)} = \liminf_{t \to \infty} t^\lambda \left[ \int_{\mathbb{R}^2} (G_t(p) - G^M(p)) |p|^4 \, dp \right] - CR^{-2} t^{\lambda'} \geq \liminf_{t \to \infty} t^\lambda \left[ \frac{1}{2} \int_{\mathbb{R}^2} (G_t(p) - G^M(p)) |p|^4 \, dp \right] > 0,
\]

where the penultimate estimate is due to (57) for sufficiently large \( t \).

**3.5. Summary of the proof of Theorem 5.** This completes the proof. The statements on the regularity of \( G \) follow from Proposition 7. The properties of the rescaling operator \( \eta \) are given in Proposition 8. The convergence was shown at the end of subsection 3.3. The lower estimate (9) was given in subsection 3.4.

**4. The Boltzmann case.** Now we consider the case where \( f \) satisfies the Boltzmann equation

\[
\partial_t f + \nabla_z f \cdot w = Q[f] \text{ for } z, w \in \mathbb{R}^2,
\]

together with the \( S \)-objectivity condition (1). For simplicity the collision operator \( Q \) is assumed to be the hard-sphere kernel

\[
Q[f](v) = \int_{S^1} \int_{\mathbb{R}^2} (f_s f'_s - f f') (v - v') \cdot \nu_+ \, dv' \, d\nu.
\]

We repeat the reduction steps in section 3 and obtain the equivalent of (12):

\[
\begin{aligned}
\partial_t g = \mu \nabla \cdot (g \alpha \otimes \beta w) + Q[g], \\
g|_{t=0} = g_0,
\end{aligned}
\]

where \( g_t = g_t(w) \) and \( Q \) is unchanged (acts on \( w \)). As before we define the kinetic energy by

\[
\theta[g] = \frac{1}{2} \int_{\mathbb{R}^2} |w|^2 g(w) \, dw.
\]

The energy \( \theta \) is conserved if and only if \( \mu = 0 \). A quantitative version of this observation delivers the existence and uniqueness of solutions for all time. For some time-dependent transformation \( \eta_t \in \mathbb{R}^{d \times d} \) we repeat the notation (4) and define

\[
\begin{aligned}
p = \eta_t w, \\
G(p) = \det \eta_t (\eta_t^{-1} p).
\end{aligned}
\]

Then the collision operator \( Q \) can be written in terms of a rescaled collision operator

\[
Q_{\eta_t}[G] = \det \eta_t Q[g],
\]
where

\[ Q_{\eta_t}[G] = \int_{S_t} \int_{\mathbb{R}^2} \int G_s G' - G G' \ \nu \cdot \eta_t^{-1}(p - p') \ \nu \cdot \eta_t^{-1}(p - p') \ d\nu \ dp' \ dv, \]
\[ p_s = p - \eta_t \nu \otimes \nu \eta_t^{-1}(p - p'), \]
\[ p'_s = p' + \eta_t \nu \otimes \nu \eta_t^{-1}(p - p'). \]

The function \( G \) satisfies the equation

\[ \partial_t G = Q_{\eta_t}[G] - \nabla_p \cdot (GFp) \]

with \( F = (\eta_t - \mu \eta_t \alpha \otimes \beta)\eta_t^{-1} \) as before.

Our results for the Boltzmann case are less detailed than for the Fokker–Planck case. Although it is not known whether (60) admits a stationary solution we can demonstrate that there is no stationary solution of exponential type.

**Theorem 13.** Equation (60) admits a global solutions if \( G_0 \in L^1 \), which preserve mass and the renormalized Cauchy stress tensor

\[ T_{\eta_t} = \frac{1}{2} \int_{\mathbb{R}^2} p \otimes p G \ dp. \]

The collision invariants of \( Q_{\eta_t} \) are 1, \( \eta_t^{-1} p \), and \( |\eta_t^{-1} p|^2 \). The solutions \( G_t \) do not converge to a function of exponential form \( K \exp(h(p)) \) with \( h(\alpha p) = \alpha h(p) \ \forall \alpha > 0, \forall p \in \mathbb{R}^2 \), and a fixed \( r > 0 \) as \( t \to \infty \). There exists a \( G_\infty \in L^1 \) with \( \|G_\infty\|_{L^1} = 1 \) and a sequence \( t_j \to \infty \) as \( j \to \infty \) such that \( G_{t_j} \) converges weakly in \( L^1 \) to \( G_\infty \).

The proof involves several parts. The collision invariants are determined in subsection 4.1, the rescaling \( \eta_t \) is determined in subsection 4.2. The shape of universal equilibria are discussed in subsection 4.3. The global bounds leading to tightness are given in subsection 4.4, which completes the proof.

**Remark 14.** The question whether (60) admits a Lyapunov functional appears to be open. It is not hard to see, if

\[ S[G] = \int_{\mathbb{R}^2} G \log G \ dp, \]

that we have

\[ \frac{dS}{dt} \leq - \frac{1}{4} \int_{\mathbb{R}^2} \int_{S^1} \min_{s \in [0, 1]} \frac{(GG' - G_s G'_s)^2}{GG' + (1 - s)GG'_s} \ [\nu \cdot \eta_t^{-1}(p - p')]_+ \ dv \ dp' \ dp - \text{tr} F. \]

However the behavior of \( F \), which will be linked to the stress rates \( P \) in (66) below, cannot be determined. Note that the first term in (61) is analogous to the standard entropy production in the case where \( \mu = 0 \). In particular, it is nonpositive. However \( \text{tr} P \) is not necessarily negative and in contrast to Remark 6 we cannot conclude that \( S \) is Lyapunov functional.
Thus, and only if \( g \in Q(62) \) is invariant. Let \( k \in \mathbb{R}^2 \) such that \( \int_{\mathbb{R}^2} (G_{\ast} G'_{\ast} - G G') \log G \left[ \nu \cdot \eta_{t}^{-1}(p - p') \right] + d\nu \, dp' \, dp - \text{tr} \, F \)
\begin{align*}
&= \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{R}^d} (G_{\ast} G'_{\ast} - G G') (\log G + \log G' - \log G_{\ast} - \log G'_{\ast}) \\
&\quad \cdot [\nu \cdot \eta_{t}^{-1}(p - p') + d\nu \, dp' \, dp - \text{tr} \, F] \\
&= \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{R}^d} (G_{\ast} G'_{\ast} - G G') (\log G' - \log G_{\ast} G'_{\ast}) \\
&\quad \cdot [\nu \cdot \eta_{t}^{-1}(p - p') + d\nu \, dp' \, dp - \text{tr} \, F] \\
&\leq \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{R}^d} \min_{s \in [0,1]} \frac{(G G' - G_{\ast} G'_{\ast})^2}{s G G' + (1 - s) G_{\ast} G'_{\ast}} [\nu \cdot \eta_{t}^{-1}(p - p') + d\nu \, dp' \, dp - \text{tr} \, F],
\end{align*}
as required. 

4.1. Collision invariants and stationary solutions. If we ignore \( \text{tr} \, F \) in (61) it is well known that the numerator vanishes if \( G \) depends only on collision invariants \( k_{\eta_t}(p) \) which are characterized by
\[
k + k' = k_{\ast} + k'_{\ast}.
\]
We determine the collision invariants below, but \( |p|^2 \) is not a collision invariant for general \( \eta_t \). Note that every collision invariant \( k \) generates a stationary solution \( G = \exp(k) \) for \( Q_{\eta_t} \), but not, in general, for the full equation (60).

**Lemma 15.** If \( G \) is a zero of \( Q_{\eta_t} \), i.e., \( Q_{\eta_t}(G) = 0 \), then there exists \( a, c \in \mathbb{R}, b \in \mathbb{R}^2 \) such that
\[
G(p) = \exp \left( a + b \cdot \eta_{t}^{-1} p + c|\eta_{t}^{-1} p|^2 \right).
\]
Let \( k_{\eta_t}(p) = p \cdot K_{\eta_t} p \) with \( K_{\eta_t} = (\eta_{t}^{\ast})^{-1} \eta_{t}^{-1} \in \mathbb{R}^{2 \times 2} \). Then \( k \) is a quadratic collision invariant.

**Proof.** Recall that
\[
(62) \quad Q_{\eta_t}(G)(p) = Q[g](\eta_{t}^{-1} p),
\]
where \( G(p) = \det \eta_{t}^{-1} g(\eta_{t}^{-1} p) \). It is well known (e.g., [8, sect. 3.2]) that \( Q[g] = 0 \) if and only if \( g = \exp(k(p)) \), where \( k \) is a collision invariant, i.e.,
\[
k(p) = a + b \cdot w + c|w|^2.
\]
Thus, \( Q_{\eta_t}(G) = 0 \) if and only if \( G = a + b \cdot \eta_{t}^{-1} p + c|\eta_{t}^{-1} p|^2 \), which is the claim. \( \square \)
4.2. Choice of rescaling. It is not hard to see that collisions do not conserve the standard kinetic energy of \( p \) and \( p' \). Indeed

\[
\begin{align*}
|p| \cdot |p'| = |p - \eta \cdot \nu \otimes \nu \eta^{-1}_{t}(p - p')| + |p' + \eta \cdot \nu \otimes \nu \eta^{-1}_{t}(p - p')| \\
= |p|^2 + |p'|^2 - 2(p - p') \cdot \eta \cdot \nu \otimes \nu \eta^{-1}_{t}(p - p') + 2(p - p') \cdot \eta^{-1}_{t} \nu \otimes \nu \eta^{2}_{t} \nu \otimes \nu \eta^{-1}_{t}(p - p')
\end{align*}
\]

(63)

so that

\[
\begin{align*}
P_{\eta} &= \frac{1}{2} \int_{\mathbb{R}^d} p \otimes p Q_{\eta} \, dp,
\end{align*}
\]

so that

\[
\begin{align*}
\text{tr} \, P_{\eta} &= \frac{1}{2} \int_{\mathbb{R}^d} \int_{\mathbb{S}^d} |p|^2 (G_{s} G_{s}^{*} - GG') \nu \cdot \eta^{-1}_{t}(p - p') \, dp \, dp' \, dp \\
&= \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{S}^d} (|p|^2 + |p'|^2) (G_{s} G_{s}^{*} - GG') \nu \cdot \eta^{-1}_{t}(p - p') \, dp \, dp' \, dp \\
&= \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{S}^d} (|p|^2 + |p'|^2) GG' \nu \cdot \eta^{-1}_{t}(p - p') \, dp \, dp' \, dp \\
&= \frac{1}{4} \int_{\mathbb{R}^d} \int_{\mathbb{S}^d} (p - p') \cdot C_{\nu} (p - p') GG' \nu \cdot \eta^{-1}_{t}(p - p') \, dp \, dp' \, dp
\end{align*}
\]

The first equation holds because the order of integration can be exchanged, the last equation follows from (63). In particular one finds that

\[
\text{tr} \, P_{\text{Id}} = 0.
\]

Our aim is to construct a time-dependent transformation \( \eta_{t} \in \mathbb{R}^{d \times d}_{\text{sym}} \) such that the renormalized Cauchy stress tensor

\[
T_{\eta} = \frac{1}{2} \int_{\mathbb{R}^d} p \otimes p G(p) \, dp
\]

is constant. We have already seen in section 3 that \( T_{\eta} = \text{Id} \) if

\[
\eta^{-2}_{t} = \frac{1}{2} \int \omega \otimes w \, g_{t}(w) \, dw,
\]

and \( g_{t} \) is a solution of (59). Differentiating the Cauchy stress with respect to \( t \) and using \( T_{\eta} = \frac{1}{2} \text{Id} \) gives

\[
\frac{dT_{\eta}}{dt} = \frac{1}{2} \int_{\mathbb{R}^d} p \otimes p (Q_{\eta} - \nabla \cdot (GFp)) \, dp = P + \frac{1}{2} (F + F^{*}) = P + F_{\text{sym}}.
\]

Thus, we have obtained a nonautonomous system of ordinary differential equations

\[
P + F_{\text{sym}} = 0.
\]

Obviously (66) is the analogue of (36).
4.3. Are there stationary solutions that are of exponential form? Assume that \( G \) is of exponential form, i.e., there exist \( h \in C^1(\mathbb{R}^2 \setminus \{0\}) \) and homogeneity exponent \( r > 0 \) so that

\[
G^{(h)}(p) = K \exp(h(p)) \quad \text{with} \quad h(\lambda p) = \lambda^r h(p) \quad \forall \lambda > 0, \forall p \in \mathbb{R}^2.
\]

Note that by integrability of \( G \) this implies \( h < 0 \) on \( \mathbb{R}^2 \). Now \( G \) is substituted into the differential equation (60). Note first that

\[
(67) \quad \nabla_p \cdot (G^{(h)}(p) \cdot F) = (\text{tr} F + \nabla h \cdot F) G^{(h)}(p) = k_i(p) G(p),
\]

where \( k \) is the sum of a constant and homogeneous function of order \( r \) in \( p \).

Furthermore in the collision term we denote

\[
p_* = p - \eta h \otimes \nu h^{-1}(p - p'), \quad p'_* = p' + \eta h \otimes \nu h^{-1}(p - p');
\]

then it has the form

\[
Q_{t_1}(G)(p) = K^2 \int_{\mathbb{R}^2} \int_{S^1} \begin{split}
\{ \exp(h(p_*)) + h(p'_*) \} - \exp(h(p) + h(p')) \left[ \nu \cdot h^{-1}(p - p') \right]_+ \, d\nu \, dp'
\end{split}
\]

\[
= K^2 \exp(h(p)) \int_{\mathbb{R}^2} \int_{S^1} \begin{split}
[\exp(h(p_*)) + h(p'_*) - h(p) - h(p')] \exp(h(p')) \left[ \nu \cdot h^{-1}(p - p') \right]_+ \, d\nu \, dp'
\end{split}
\]

\[
= K^2 \exp(h(p)) \int_{\mathbb{R}^2} \exp(h(p - q)) \int_{S^1} \begin{split}
\exp(h(p - \nu) \otimes \nu h^{-1}(q)) + h(p - q + \nu) \otimes \nu h^{-1}(q) - h(p) - h(p - q) - 1 \end{split}
\]

\[
\left[ \nu \cdot h^{-1}(q) \right]_+ \, d\nu \, dq.
\]

To cancel the expression in (67), the integrals over \( j \) are necessarily \( O(|p|^r) \) as \( |p| \to \infty \). First we consider the case when \( q = p \). Then the exponent can be simplified to

\[
i(\nu) := h(p - \eta \nu \otimes \nu h^{-1}(p) + h(\eta \nu \otimes \nu h^{-1}(p) - h(p) - h(0)).
\]

If there are \( \nu \) and \( p \) such that

\[
(68) \quad i(\nu) > 0,
\]

then \( j(q) \) grows exponentially for a sector of \( q \) in \( \mathbb{R}^2 \). Due to continuous dependence of this sector on \( \nu \), the integral \( \int_{S^1} j(q) \, d\nu \) still grows exponentially on some sector in \( \mathbb{R}^2 \). By choosing \( p \) in such a sector, we obtain constants \( c_1, c_2 \) such that

\[
Q_{t_1}(G)(p) \geq (c_1 j(p) - c_2) G(p),
\]

which cannot be equal to a term \( k_i(p) G(p) \) as in (67), where \( k \) is bounded by a polynomial.

We now establish the existence of some \( \nu \) such that (68) holds. For \( t = 0 \) we have \( \eta_t = \text{Id} \) and by (65) and (66) \( \text{tr} F = 0 \); then (60) and (67) imply for \( p = 0 \) that

\[
0 = Q_{t_1}(G^{(h)}) = Q[G^{(h)}]
\]

\[
= K^2 \int_{\mathbb{R}^2} \exp(-q) \int_{S^1} \begin{split}
(\exp(-\nu \otimes \nu q) + h(-q + \nu \otimes \nu q) - h(-q)) - 1 \end{split}
\]

\[
\left[ \nu \cdot q \right]_+ \, d\nu \, dq.
\]

Unless \( r = 2 \), when \( h \) is a collision invariant, this implies that the exponent attains positive and negative values. By exchanging the roles of \( p \) and \( p' \), we hence obtain that
(68) will hold for some \( p \) and \( \nu \), hence ruling out any \( h \) with homogeneity exponent \( r > 0, r \neq 2 \).

Now consider the only remaining case \( h(p) = -d|\eta|^{-1}p^2 \) for some \( d > 0 \). The collision invariance of \( h \) implies \( Q_{\eta_t}[G^{(h)}] \equiv 0 \). Furthermore \( \eta_t = \text{Id} \) is constant as \( G^{(h)} \) is constant, i.e., \( \partial_t G^{(h)} = 0 \). Hence we obtain the equation

\[
\nabla_p \cdot (G^{(h)}(p) \cdot Fp) = (\text{tr} F + 2dp \cdot Fp)G^{(h)} = 0.
\]

For \( \eta_t = \text{Id} \), \( F = -\mu \alpha \otimes \beta \), which is incompatible with \( \text{tr} F - 2dp \cdot Fp = -2dp Fp = 0 \), thus ruling out the collision invariant.

**4.4. Completing the proof of Theorem 13.** The regularity of the solution follows from the regularity proposition below.

**Proposition 16.** If \( \int_{\mathbb{R}^2} g_0(w)(1 + |w|^2) \, dw < \infty \), then (59) admits a unique mild solution for all \( t > 0 \).

**Proof.** The transport term \( \mu v \partial_v g \) is the generator of the strongly continuous semigroup \( X_t : L^1(\mathbb{R}^2) \to L^1(\mathbb{R}^2) \) on the space of integrable function that satisfy \( \int_{\mathbb{R}^2} g_0(w)(1 + |w|^2) \, dw < \infty \). The semigroup is given explicitly by \( X_t g(w) = g((\text{Id} + \mu t \alpha \otimes \beta)w) \). Furthermore by Povzner’s inequality \( Q \) is a continuous nonlinear operator on \( L^2(\mathbb{R}^2) \). This gives the existence of the unique mild solution to (59) given by

\[
g_t = X_t g_0 + \int_0^t X_{t-s} Q(g_s, g_s) \, ds.
\]

We show that we can continue this solution globally by showing that

\[
\int_{\mathbb{R}^2} g_t(w)(1 + |w|^2) \, dw < \infty
\]

for all times. First recall that

\[
\int_{\mathbb{R}^2} |w|^2 Q[g] \, dw = 0
\]

holds for any density \( g \) with \( \theta[g] = \frac{1}{2} \int |w|^2 \, dw < \infty \) because \( |w|^2 \) is a collision invariant. Differentiating \( \theta \) with respect to \( t \) gives

\[
\frac{d\theta}{dt} = \frac{1}{2} \int_{\mathbb{R}^2} |w|^2 (\mu \nabla \cdot (g \alpha \otimes \beta w) + Q[g]) \, dw
\]

\[
= \frac{1}{2} \int_{\mathbb{R}^2} |w|^2 Q[g] \, dw - \mu \int_{\mathbb{R}^2} g \, w \cdot \alpha \otimes \beta w \, dw
\]

\[
\leq \frac{\mu}{2} \int_{\mathbb{R}^2} |w|^2 g \, dw = |\mu| \theta[g].
\]

Thus, \( \theta[g] \leq e^{\mu |t|} \theta[g_0] \). By a similar argument, \( \int_{\mathbb{R}^2} g_t(w) \, dw = \int_{\mathbb{R}^2} g_0(w) \, dw \), so \( \int_{\mathbb{R}^2} g_t(w)(1 + |w|^2) \, dw \) remains bounded for bounded times, such that (69) defines a global mild solution, which is unique by a Gronwall argument.

Then we transform the mild solution \( g \) as in the paragraph preceding (60) to obtain a global solution \( G \) of (60).

The choice of \( \eta_t \) in subsection 4.2 gives the preservation of mass and energy for \( G_t \); this immediately gives the weak convergence of subsequences to some limit
points. The collision invariants are characterized in Lemma 15. The shape of a possible equilibrium is analyzed in subsection 4.3.

5. Conclusion. We studied two closely related equations in kinetic theory, the Fokker–Planck equation and the Boltzmann equation with shear boundary conditions. The boundary conditions are not compatible with the conservation of energy. After rescaling the velocities in an anisotropic fashion we obtain renormalized equations which have the property that solutions conserve all second moments and, in particular, the energy. The renormalized Fokker–Planck equation admits Maxwellian equilibria and the long-time behavior of renormalized solutions can be characterized completely. More precisely, we show rigorously that as $t$ tends to infinity solutions converge at an algebraic rate to the Maxwellian with the appropriate second moments.

On the other hand the renormalized Boltzmann equation does not admit equilibria of exponential type including Maxweillians. Indeed, due to the nonautonomous nature of the shape equation (60) there might be no equilibria at all. We conjecture that for large time, solutions of the renormalized Boltzmann converge to a limiting density, but a rigorous proof is not available.

Results on the existence of self-similar profiles (i.e., equilibria for the shape equation) and long-time behavior in the case of soft interaction potentials have been obtained in [23] and [24]. In [23] the existence of stationary self-similar solutions is established rigorously for Maxwellian molecules (where the repulsive force between particles at distance $r$ is $r^{-5}$) after isotropic rescaling (where $\eta$ is a multiple of the identity). Detailed information about energy flux can then be derived. Moreover, in [24] formal calculations covering the supercritical case where the force decays faster than $r^{-5}$ are being presented. Based on these calculations the authors conjecture that after isotropic rescaling in the supercritical case solutions converge to a Maxwellian.

It is noteworthy that the analysis in [23] and [24] also covers other objectivity conditions than the simple shear, for example, homogeneous dilations where $S = (\text{Id}, -\text{Id})$, and other choices. In view of these results it will be worthwhile to extend our approach to other objectivity conditions and explore different choices for the collision operator in the Boltzmann equation.


Proof of Proposition 4. Assume that $h(w) = a + b \cdot w$ is affine. Observe that (1) implies for each $w$ that

\[
\nabla_z f(x, w) \cdot \alpha = 0, \\
\nabla_z f \cdot \beta = -\mu \nabla_w f \cdot \alpha.
\]

Consider now the quantity

\[
H = \int_{\mathbb{R}^2} h(w) \partial_t f(x, w) \, dw \\
= \int_{\mathbb{R}^2} h(w) \left( \Delta_w f + \rho \theta^{-1} \nabla_w \cdot (f(x, w)(w - \rho^{-1} \dot{v})) - \nabla_z f(x, w) \cdot w \right) \, dw \\
= \int_{\mathbb{R}^2} \left\{ \left( \Delta_w h - \rho \theta^{-1} \nabla h \cdot (w - \rho^{-1} \dot{v}) \right) f(x, w) + \mu \left( \nabla_w f \cdot \alpha \right)(w \cdot \beta) \right\} \, dw.
\]

Clearly $\Delta h = 0$ as $h$ is affine. Moreover $\int_{\mathbb{R}^2} b \cdot (w - \rho^{-1} \dot{v}) f(w) \, dw = 0$ by the definition of $\dot{v}$ and finally $\int_{\mathbb{R}^2} \left( \nabla_w f(x, w) \cdot \alpha \right)(w \cdot \beta) \, dw = 0$ by partial integration.
This implies that
\[ t \mapsto \int_{\mathbb{R}^2} h(w) f_t(x, w) \, dw \text{ is constant} \]
for all \( x \in \mathbb{R} \) and by (1) this is constant and thereby the first claim.

Next, assume that \( \mu = 0 \) and \( h(w) = \frac{1}{2} |w|^2 \). Repeating the previous calculation we obtain
\[
H = \int_{\mathbb{R}^2} h(w) \partial_t f(x, \alpha, w) \, dw \]
\[ = \int_{\mathbb{R}^2} (2 - \theta^{-1} |w|^2) f(x, \alpha, w) \, dw = \int_0^1 (2 \rho(x, \alpha) - 2 \theta^{-1} \rho \theta(x, \alpha)) \, dw = 0. \]

Finally we demonstrate that \( f^M \) is a stationary solution. One finds that
\[
\rho = -\frac{\pi}{c} \exp \left( a - \frac{|b|^2}{4c} \right), \quad w_0 = \frac{\pi}{2c^2} \exp \left( a - \frac{|b|^2}{4c} \right) b, \quad \theta = \frac{\pi}{2c^2} \exp \left( a - \frac{|b|^2}{4c} \right); \]
in particular, \( \rho \theta^{-1} = -2c \) and \( \rho^{-1} w_0 = -\frac{1}{2c} b \). Then
\[
Lf^M = \Delta w f^M + \rho \theta^{-1} \nabla_w \cdot (f^M (w - \rho^{-1} w_0)) - \nabla_z f^M \cdot w \\
= \left( |\nabla h|^2 + \Delta h - 4c - 2c \nabla h(w) \cdot \left( w + \frac{1}{2c} b \right) \right) f^M \\
= \left( 4c^2 |w|^2 + |b|^2 + 4c b \cdot w + 4c - 4c(2c w + b) \cdot \left( w + \frac{1}{2c} b \right) \right) f^M \\
= \left( (4c^2 - 4c^2) |w|^2 + (4c b - 2c (b + b)) \cdot w + |b|^2 - |b|^2 \right) f^M = 0.
\]

To prove convergence for a general spatially homogeneous initial datum \( f_0 \), we rewrite the equation in a spirit similar to subsection 3.3 and (38). Using that mass \( \rho \), momentum \( w_0 \), and energy \( \theta \) remain constant along solutions, choose \( f^M \) such that its triple \( \rho, w_0, \theta \) coincides with the triple of \( f_0 \). Then write \( f = u f^M \) with \( u \in L^2(\mathbb{R}^2, df^M) \). To show \( L^1 \) convergence it is enough to show that \( u \to 0 \) in \( L^2(\mathbb{R}^2, df^M)/\text{span}(1) \) by Hölder’s inequality. The relative profile \( u \) satisfies the equation
\[
\partial_t u = \Delta u + (b + 2cw) \cdot \nabla u = -A^* Au,
\]
where \( A = \nabla u \) and \( A^* = -\nabla \cdot (b + 2cw) \) is its adjoint operator in \( L^2(\mathbb{R}^2, df^M)/\text{span}(1) \) with inner product \( \langle \cdot, \cdot \rangle \). Then
\[
\partial_t \langle u, u \rangle = 2\langle \partial_t u, u \rangle = -2 \langle Au, Au \rangle \leq -2C \langle u, u \rangle
\]
using a Poincaré inequality as in [33, A.19], which then gives the required exponential convergence of \( u \) to 0 in \( L^2(\mathbb{R}^2, df^M)/\text{span}(1) \).

**Proof of Remark 6.** We calculate \( \frac{d}{dt} S[G] \) along solutions of (7), noticing that \( G \) is smooth with respect to \( p \) as \( g \) is smooth for \( t > 0 \), so we can perform integration.
by parts (ibp), etc.,

\[
\frac{d}{dt} S[G_t] = \frac{d}{dt} \int_{\mathbb{R}^2} G_t(p) \ln \frac{G_t(p)}{\exp(-|p|^2/2)} \, dp = \int_{\mathbb{R}^2} \left( \ln \frac{G_t(p)}{\exp(-|p|^2/2)} + 1 \right) \partial_t G_t(p) \, dp
\]

\[
= \int_{\mathbb{R}^2} \left( \ln \frac{G_t(p)}{\exp(-|p|^2/2)} + 1 \right) \nabla_p \cdot (G_t(p) (\theta^{-1} \text{Id} - F) p + T^{-1} \nabla G_t(p)) \, dp
\]

\[
= - \int_{\mathbb{R}^2} \nabla \left( \ln \frac{G_t(p)}{\exp(-|p|^2/2)} + 1 \right) \cdot (G_t(p) (\theta^{-1} \text{Id} - F) p + T^{-1} \nabla G_t(p)) \, dp
\]

\[
= - \int_{\mathbb{R}^2} G_t(p) \left( \nabla G_t(p) + G_t(p)p \cdot (G_t(p) (\theta^{-1} \text{Id} - F) p + T^{-1} \nabla G_t(p)) \right) \, dp
\]

Next we split \( F \) into its symmetric and antisymmetric parts, which are given by \( \frac{1}{2}(F + F^*) \) and \( \frac{i}{2}(F - F^*) \), respectively. For the symmetric part we use (36) and find that

\[
= - \int_{\mathbb{R}^2} G_t(p) \left( \nabla G_t(p) + G_t(p)p \cdot (G_t(p) (\theta^{-1} \text{Id} - F) p + T^{-1} \nabla G_t(p)) \right) \, dp
\]

\[
= - \int_{\mathbb{R}^2} G_t(p) \left( \nabla G_t(p) + G_t(p)p \cdot T^{-1} (G_t(p)p + \nabla G_t(p)) \right) \, dp
\]

\[
+ \frac{1}{2} \int_{\mathbb{R}^2} G_t(p) \left( \nabla G_t(p) + G_t(p)p \cdot G_t(p) (F - F^*)p \right) \, dp
\]

\[
= - \int_{\mathbb{R}^2} G_t(p) \left| \eta_t (G_t(p)p + \nabla G_t(p)) \right|^2 \, dp + \frac{1}{2} \int_{\mathbb{R}^2} \nabla G(p) \cdot (F - F^*)p \, dp
\]

\[
+ \frac{1}{2} \int_{\mathbb{R}^2} p \cdot (F - F^*)p G_t(p) \, dp
\]

\[
= - \int_{\mathbb{R}^2} G_t(p) \left| \eta_t (G_t(p)p + \nabla G_t(p)) \right|^2 \, dp;
\]

the other two integrals are zero, the middle one by integration by parts and the final one due to the antisymmetry of \( F - F^* \).

Hence \( S[G_t] \) decays unless \( G_t(p)p + \nabla G_t(p) = 0 \); the only differentiable solutions in \( L^1(\mathbb{R}^2) \) are multiples of the Maxwellian \( G^M \). \( \square \)

REFERENCES


