Slowly Varying Modes in a 2D Duct with Shear Flow and Lined Walls

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A slowly varying modes solution of WKB type has been derived for the problem of sound propagation in a slowly varying 2D duct with homentropic inviscid sheared mean flow and acoustically lined walls. The slowly varying modal shape function and axial wave number are described by the Pridmore-Brown eigenvalue equation.

The slowly varying modal amplitude was determined in the usual way by an equation resulting from a solvability condition. For a general mean flow, this equation could be solved in the form of an incomplete adiabatic invariant. Due to conservation of specific mean vorticity along streamlines, two simplifications proved possible for a linearly sheared mean flow: (1) an analytically exact approximation for the mean flow could be found, and (2) the solvability condition reduced to a complete adiabatic invariant for the acoustics.

For this last configuration some example cases were evaluated numerically, where the Pridmore-Brown eigenvalue problem was solved by a Galerkin projection combined with an efficient nonlinear iteration.

Key words: TBD

1. Introduction

Sound propagation in uniformly lined straight ducts with uniform mean flow is well established by its analytically exact description in duct modes (Rienstra 2003a, 2015; Rienstra & Hirschberg 1992). In cross-wise nonuniform flow there are still modes, although the (Pridmore-Brown) equation that describes them is in general not solvable in terms of standard functions (Pridmore-Brown 1958; Vilenski & Rienstra 2007a,b; Brambley, Rienstra & Darau 2012; Oppeneer et al. 2011; Oppeneer, Rienstra & Sijtsma 2016; Rienstra 2019). An exception is 2D linear shear flow and uniform soundspeed, for which it can be solved exactly by Weber’s Parabolic Cylinder functions (Goldstein & Rice 1973).

These modal solutions provide insight, but the important effects due to the variation of the duct geometry and the corresponding mean flow cannot be described exactly. At least in slowly varying ducts a systematic approximate solution is possible. The inherently smooth variation of a flow duct provides a small parameter $\varepsilon$ (the slenderness of the duct wall variation) that allows a multiple scales approach for mean flow and slowly varying modes of WKB (Wentzel-Kramers-Brillouin) type (Nayfeh & Telionis 1973; Rienstra 1999, 2000; Cooper & Peake 2001; Peake & Cooper 2001; Ovenden 2002; Rienstra 2003b; Ovenden et al. 2004; Ovenden & Rienstra 2004; Ovenden 2005; Brambley & Peake 2008; Lloyd & Peake 2013). In the WKB approach it is assumed that the wave retains

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its shape. If the medium changes slowly enough (meaning little change over one wave length) the mode will not jump to another eigenstate. The wave function will gradually stretch or compress but the number of nodes and anti-nodes will not change. Under suitable conditions this may be represented mathematically by conserved quantities called adiabatic invariants. They may be related to conserved energy, but not necessarily. If an adiabatic invariant can be found (this may be far from trivial) the WKB solution is almost as transparent as its straight duct counterpart.

The procedure for slowly varying modes in mean flow is, broadly speaking, as follows. First the mean flow is assumed to be slowly varying everywhere (no entrance effects) so the mean flow equations are rewritten in the slow coordinate $X = \varepsilon x$, rescaled and simplified to leading order in $\varepsilon$. These equations are solved analytically if possible (apart from an algebraic equation, this is usually the case for potential flow: Rienstra (1999); Peake & Cooper (2001); Rienstra (2003b); Brambley & Peake (2008)), or otherwise numerically (often for rotational flow: Cooper & Peake (2001); Lloyd & Peake (2013)). Then for the acoustic equations a slowly varying mode is assumed† of the form $N(X)\psi(x, y, z)\exp(-i\int \kappa(\varepsilon z)dz)$, where $\psi$ is conveniently normalised. Its existence requires in WKB-theory an equation (also known as solvability condition) of the form

$$\frac{d}{dX} (N^2 F) = N^2 G$$

where $F$ and $G$ are functions of $X$ that depend on cross-sectional integrals of squares and products of $\psi, \psi_X, \psi_y$, and $\psi_z$. For second order linear equations this is always possible. If a form can be found without the presence of $\psi_X$, the relation between $N$ and $\psi$ is really a local property and we call this an incomplete adiabatic invariant, as the solution for $N^2$ includes an integration of $G/F$ to $X$. If also $G = 0$, the solvability condition is completely integrable and we call the resulting integral $N^2 F = \text{constant}$ a complete adiabatic invariant. In that case the scaling of $\psi$ may be the invariant itself.

After the solutions for irrotational flow (Rienstra 1999; Peake & Cooper 2001; Rienstra 2003b; Brambley & Peake 2008), great progress was made by Peake et al. for rotational flow in cylindrical ducts, with swirl (Cooper & Peake 2001) and without swirl (Lloyd & Peake 2013). The complexity of the 3D geometry and mean flows, however, prevents both the existence of an adiabatic invariance and an analytically exact mean flow.

In the present paper we aim to extend Peake and colleagues’ theory for slowly varying modes in non-uniform shear flow. We will show that it is possible to make progress by considering a 2D duct. For the general situation an incomplete adiabatic invariant could be found, while for the particular case of a linearly sheared mean flow the mean flow equations can be solved analytically exactly and for the acoustics a complete adiabatic invariant exists.

The present work builds on, but at the same time improves and replaces, a previous WKB analysis of the related problem with hard walls, published as conference paper (Rienstra 2016). Apart from some minor notational differences (like $\Omega$ for $\Omega/C$, $\Psi$ for $M$, $\kappa$ for $\mu$) and algebraic corrections, the main differences are the lined walls, the found adiabatic invariants, and the use of a numerical solution of the central Pridmore-Brown equation, giving access to a very wide range of parameters.

† Equivalent to suppression of secular terms.
2. The problem

2.1. The equations

In the acoustic, i.e. isentropic realm of a perfect gas that we will consider, we have for pressure $\tilde{p}$, velocity $\tilde{v}$, density $\tilde{\rho}$, entropy $\tilde{s}$, and soundspeed $\tilde{c}$

\[
\begin{align*}
\frac{d\tilde{\rho}}{dt} &= -\tilde{\rho} \nabla \cdot \tilde{v}, \quad \frac{d\tilde{v}}{dt} = \nabla \tilde{p}, \\
\frac{d\tilde{s}}{dt} &= C_V \frac{d\tilde{p}}{\tilde{\rho} dt} - C_P \frac{d\tilde{p}}{\tilde{\rho} dt} = 0, \quad \tilde{c}^2 = \frac{\gamma \tilde{p}}{\tilde{\rho}}, \quad \gamma = \frac{C_P}{C_V},
\end{align*}
\]

(2.1)

where $\gamma$, $C_P$, and $C_V$ are gas constants. $C_V$ is the heat capacity at constant volume, $C_P$ is the heat capacity at constant pressure, and $\gamma = C_P/C_V$. With isentropy, entropy $\tilde{s}$ is constant along a streamline. We will assume a bit more, namely the constant $\tilde{s} = s_{ref}$ is the same for all streamlines, which is homentropic flow. In that case we can integrate the entropy equation to get

\[
C_V \log \tilde{p} - C_P \log \tilde{\rho} = s_{ref} = C_V \log p_{ref} - C_P \log \rho_{ref}
\]

or

\[
\frac{\tilde{p}}{p_{ref}} = \left( \frac{\tilde{\rho}}{\rho_{ref}} \right)^{\gamma}.
\]

(2.2)

When we have a stationary mean flow with unsteady time-harmonic perturbations of frequency $\omega$, given, in the usual complex notation, by

\[
\tilde{v} = V + \text{Re}(v e^{i\omega t}), \quad \tilde{p} = P + \text{Re}(p e^{i\omega t}), \quad \tilde{\rho} = D + \text{Re}(\rho e^{i\omega t}),
\]

(2.3)

and linearise for small amplitude, we obtain for the mean flow

\[
\nabla \cdot (DV) = 0, \quad D(V \cdot \nabla)V = -\nabla P,
\]

\[
\frac{P}{p_{ref}} = \left( \frac{D}{\rho_{ref}} \right)^{\gamma}, \quad C^2 = \frac{\gamma P}{D}.
\]

(2.4)

The mean flow momentum equation can be integrated along streamlines to yield (Batchelor 1967) a version of Bernoulli’s equation that the total specific enthalpy

\[
\frac{1}{2} |V|^2 + \frac{C^2}{\gamma - 1} = \text{constant along a streamline},
\]

(2.5)

not necessarily the same constant for all streamlines. Another quantity, conserved along streamlines, that will be of interest is the specific vorticity $\Xi/D$, where $\Xi = V_x - U_y$. Take the curl of the mean flow momentum equation and substitute the mass conservation to get, by virtue of the pressure being barotropic (Batchelor 1967),

\[
V \cdot \nabla \left( \frac{\Xi}{D} \right) = 0.
\]

(2.6)

Finally, the perturbations are described by the equations

\[
\begin{align*}
&i\omega \rho + \nabla \cdot (V \rho + vD) = 0, \\
&D(i\omega + V \cdot \nabla)v + D(v \cdot \nabla)V + \rho(V \cdot \nabla)V = -\nabla p, \\
&p = C^2 \rho.
\end{align*}
\]

(2.7)

Myer’s time-averaged acoustic energy equations for inviscid and homentropic flow, with $\Xi$ and $\xi$ the mean and perturbed flow vorticity, are given by (Myers 1991; Rienstra &
\[ y = h(\varepsilon x) \]
\[ y = g(\varepsilon x) \]

Figure 1: Sketch of geometry: 2-dimensional slowly varying duct

Hirschberg 1992)

\[ \nabla \langle I \rangle = -\langle D \rangle \]
\[ \langle I \rangle = \frac{1}{2} \Re \left[ (Dv + \rho V) \left( \frac{p^*}{D} + V \cdot v^* \right) \right] , \]
\[ \langle D \rangle = -\frac{1}{2} \Re \left[ \xi D \left( \frac{v^*}{-v^*} \right) + \Xi \rho v^* \cdot (-V, U) \right] . \]

2.2. Nondimensionalisation

Without further change of notation, we will assume throughout this paper that the problem is made dimensionless: lengths on a typical duct height \( a_{\text{ref}} \), velocities on a typical sound speed \( c_{\text{ref}} \), time on \( a_{\text{ref}} / c_{\text{ref}} \), densities on a typical density \( \rho_{\text{ref}} \), and pressures on \( \gamma p_{\text{ref}} = \rho_{\text{ref}} c_{\text{ref}}^2 \). The dimensionless equations are the same as above, except

\[ \gamma P = DC^2 = D^\gamma \]

while dimensionless frequency \( \omega \) is also known as the Helmholtz number.

2.3. The geometry

The domain of interest is a 2D duct \( \mathcal{V} \) of slowly varying cross section \( g(X) \leq y \leq h(X) \), where \( 0 < h - g = O(1) \). \( X = \varepsilon x \) is a so-called slow variable while \( \varepsilon \) is small. At the duct top surface \( y = h \) and the bottom surface \( y = g \), the gradients

\[ \nabla (y - h(X)) = e_y - \varepsilon e_x h_X, \quad \nabla (y - g(X)) = e_y - \varepsilon e_x g_X. \]

are vectors normal to the surface, so the outward normals of the top and bottom surface are given by

\[ n_h = \frac{e_y - \varepsilon e_x h_X}{\sqrt{1 + \varepsilon^2 h_X^2}} = e_y - \varepsilon e_x h_X + O(\varepsilon^2), \]
\[ n_g = \frac{-e_y - \varepsilon e_x g_X}{\sqrt{1 + \varepsilon^2 g_X^2}} = -e_y + \varepsilon e_x g_X + O(\varepsilon^2). \]

Evidently, \( e_x \) and \( e_y \) denote the unit vectors in \( x \) and \( y \) direction, the indices \( h \) and \( g \) refer to the top and bottom walls, and the index \( X \) to a partial derivative to \( X \). Occasionally, the index \( y \) will also be used to denote a partial derivative to \( y \).

2.4. Boundary conditions

If we denote the mean flow by \( V = U e_x + V e_y \), the impermeable duct wall yields the mean flow boundary condition \((V \cdot n) = 0 \), or

\[ V - \varepsilon g_X U = 0 \text{ at } y = g(X), \quad \text{and} \quad V - \varepsilon h_X U = 0 \text{ at } y = h(X). \]

(2.10)
The mean flow mass flux, given by
\[
\int_{g}^{h} DU \, dy = \mathcal{F},
\]
(2.11)
is then independent of \(x\), where \(\mathcal{F}\) is (after non-dimensionalisation) a constant \(O(1)\). The mean flow is assumed to be determined by the slowly varying geometry only.

The acoustic boundary condition of the lined wall includes the refraction through the vanishing boundary layer, and is a form of the so-called Ingard-Myers condition (Ingard 1959; Myers 1980), given at \(y = h (i = h)\) and \(y = g (i = g)\) by
\[
i\omega Z_i(v \cdot n_i) = (i\omega + V \cdot \nabla - n_i \cdot (n_i \cdot \nabla))p.
\]
(2.12)

3. Mean flow

3.1. The general case

Since we assumed the mean flow to be determined by the slowly varying geometry only, we write all flow variables as a function of \((X, y; \varepsilon)\), and expand each in a regular Poincaré expansion in powers of \(\varepsilon^2\) (the small parameter that appears in the equations). From elementary order of magnitude considerations it follows that \(U = O(1)\), \(V = O(\varepsilon)\), \(D = O(1)\), \(C = O(1)\), and \(P = O(1)\). So we have
\[
U = U_0 + O(\varepsilon^2), \quad V = \varepsilon V_0 + O(\varepsilon^3),
\]
(3.1)
\[
D = D_0 + O(\varepsilon^2), \quad P = P_0 + O(\varepsilon^2), \quad C = C_0 + O(\varepsilon^2),
\]
where each term in the expansion is independent of \(\varepsilon\). For notational convenience we leave out the 0 subscript, because higher orders will not be considered.

We substitute these expansions in the conservation equations and collect the terms of like powers of \(\varepsilon\). Then we get to leading order
\[
(DU)_X + (DV)_y = 0,
\]
\[
D(UU_X + VU_y) + P_X = 0,
\]
(3.2a)
\[
P_y = 0,
\]
with boundary conditions
\[
V - g_X U = 0 \text{ at } y = g \text{ and } V - h_X U = 0 \text{ at } y = h(X).
\]
(3.2b)

From \(P = P(X)\) is also \(D = D(X)\) and \(C = C(X)\), and so \(P_X = C^2 D_X = D^{\gamma-1} D_X\). From the flux condition (2.11) we find
\[
D(X) = \frac{\mathcal{F}}{\int_{g}^{h} U(X, y) \, dy}.
\]
(3.3)

Furthermore, the mass and momentum equations become
\[
(DU)_X + DV_y = 0,
\]
\[
DUU_X + DVU_y + D^{\gamma-1} D_X = 0,
\]
(3.4)
of which the last equation can be written as
\[
U \frac{\partial}{\partial X} \left( \frac{1}{2} U^2 + \frac{D^{\gamma-1}}{\gamma-1} \right) + V \frac{\partial}{\partial y} \left( \frac{1}{2} U^2 + \frac{D^{\gamma-1}}{\gamma-1} \right) = 0
\]
This again amounts to the approximated Bernoulli’s equation (2.5), saying that
\[ \frac{1}{2}U^2 + \frac{D^{\gamma-1}}{\gamma - 1} = \text{constant along streamlines}, \tag{3.5} \]
for example the walls \( y = g \) and \( y = h \).

By combining the mass and \( y \)-derivative of the momentum equations, we find the approximated version of (2.6)
\[ U \frac{\partial}{\partial X} \left( \frac{U_y}{D} \right) + V \frac{\partial}{\partial y} \left( \frac{U_y}{D} \right) = 0 \]
which means that the specific mean shear
\[ \frac{U_y}{D} = \text{constant along streamlines}. \tag{3.6} \]
If we start at (say) \( X = X_0 \) with a linearly sheared flow, i.e. \( U_y = \sigma_0 \), then the conserved constant \( \lambda = \sigma_0/D(X_0) \) is the same for all \( X \) and \( y \), and \( U_y = \lambda D \) everywhere.

The equation of mass conservation leads to
\[ DV = g_X DU(g) - \int_g^y (DU)_X \, dy', \tag{3.7} \]
which satisfies the boundary condition at \( y = g \) (by construction), and at \( y = h \) because
\[ DV - DU h_X = -\frac{d}{dX} \int_g^h DU \, dy = -\frac{d}{dX} F = 0. \]
For a potential mean flow the above conditions would be sufficient to define the flow (Rienstra 2003b). For example, \( U \) can only be constant in \( y \), and the other variables follow by an algebraic equation to be solved at each position \( X \). For the sheared (vortical) mean flow we have here, we need to define an initial velocity profile \( (U, \varepsilon V) \) and integrate numerically (iteratively) along streamlines while using (3.3), (3.5) and (3.7). This is not straightforward. See for example Lloyd & Peake (2013) for the closely related cylindrical version of this problem where use is made of the streamfunction.

### 3.2. Linearly sheared flow

A numerical evaluation of the mean flow is, however, not always necessary. Hinted by (3.6) and its following paragraph, it appears that for a velocity profile that is linear in \( y \), we can find a solution that avoids a numerical integration and much alike the potential flow solution requires no more than solving an algebraic equation. Even more, an accidental property will, rather unexpectedly, give our acoustic solution for this case a complete adiabatic invariant.

The type of velocity profile we will consider is to leading order given by
\[ U(X, y) = \tau(X) + \sigma(X)(y - g(X)), \tag{3.8} \]
where \( \sigma \) and \( \tau \) are to be determined from the conditions of fixed mass flux and varying wall normal. Then
\[ DV = DG_X U - (D\tau)_X (y - g) - \frac{1}{2} (D\sigma)_X (y - g)^2 \]
and
\[ D\tau \tau_X + D^{\gamma-1} D_X + (D\sigma_X - D_X \sigma) (\tau(y - g) + \frac{1}{2} \sigma(y - g)^2) = 0. \]
This requires (at \( y = g \)) that
\[
D \tau \tau_X + D^{\gamma - 1} D_X = 0,
\]
and therefore
\[
D \sigma_X - D_X \sigma = 0,
\]
so \( \sigma/D = \text{constant} \), which agrees with (3.6). Altogether we find that for any (physically allowable) constants \( \lambda \) and \( E \) a mean flow exists given by
\[
\frac{1}{2} \tau^2 + \frac{1}{\gamma - 1} D^{\gamma - 1} = E,
\]
(Bernoulli’s equation (3.5) along streamline \( y = g \)) and
\[
\sigma(X) = \lambda D(X).
\]
(3.9)

Eliminate \( \tau \) from
\[
\int_{g}^{h} DU \, dy = D \int_{g}^{h} U \, dy = D(h-g)(\tau + \frac{1}{2} \sigma(h-g)) = F
\]
to get†
\[
\tau = \frac{F}{D(h-g)} - \frac{1}{2} \sigma(h-g).
\]
For given geometry \( h \) and \( g \) and mean flow constants \( E \), \( \lambda \) and \( F \) we have
\[
\frac{1}{2} \left( \frac{F}{D(h-g)} - \frac{1}{2} \lambda D(h-g) \right)^2 + \frac{1}{\gamma - 1} D^{\gamma - 1} = E,
\]
(3.10)
to be solved per \( X \) for \( D \). Then we have \( \tau \), \( \sigma \), \( U \) and \( C \) like given above, from which \( V \) is found to be
\[
V = g_X U(y) + U(h) \frac{C^2 - \tau U(y)}{C^2 - \tau U(h)} \frac{h_X - g_X}{h - g}(y-g).
\]
(3.11)

Note that there is apparently no slowly varying mean flow of linear shear with \( \tau(X) \equiv 0 \).

4. Acoustic field

4.1. Slowly varying modes

With the mean flow variables expanded to leading order, and after eliminating \( \rho = p/C^2 \), we have the acoustic equations
\[
i \omega p + Up_x + DC^2(u_x + v_y) = -\varepsilon \left[ C^2 \left( \frac{U}{C^2} \right)_x p + C^2 D_X u + (Vp)_y \right] + O(\varepsilon^2),
\]
\[
D(i \omega u + U u_x) + DU_y v + p_x = -\varepsilon \left[ -D^{-1} D_X p + D(U_X u + U u_y) \right] + O(\varepsilon^2),
\]
\[
D(i \omega v + U v_x) + p_y = -\varepsilon D(Vv)_y + O(\varepsilon^2),
\]
(4.1a)
and the lined wall boundary condition at \( y = h(X) \), respect. at \( y = g(X) \),
\[
i \omega Z_h(v - \varepsilon h_X u) = \left( i \omega + U \frac{\partial}{\partial x} + \varepsilon V \frac{\partial}{\partial y} - \varepsilon V_y + \varepsilon h_X U_y \right)p + O(\varepsilon^2),
\]
\[
-i \omega Z_g(v - \varepsilon g_X u) = \left( i \omega + U \frac{\partial}{\partial x} + \varepsilon V \frac{\partial}{\partial y} - \varepsilon V_y + \varepsilon g_X U_y \right)p + O(\varepsilon^2).
\]
(4.1b)
† Indeed, with this \( \tau \) Bernoulli’s equation (3.5) also applies along streamline \( y = h \).
The assumption of a multiple scales solution is here equivalent to the WKB-Ansatz:

\[ [p, u, v] = \left[ \bar{\Psi}(X, y; \varepsilon), A(X, y; \varepsilon), B(X, y; \varepsilon) \right] e^{-i f^\varepsilon_{x z} dz}. \tag{4.2} \]

Introduce

\[ \Omega = \frac{\omega - \kappa U}{C}, \quad \Omega_y = -\frac{\kappa U_y}{C}, \tag{4.3} \]

and substitute (4.2) in (4.1) to obtain after some simplifications to \( O(\varepsilon^2) \)

\[ \Omega^2 \Psi - \kappa D\Omega CA - iD\Omega CB_y = i\varepsilon\Omega C \left[ \frac{U\Psi}{C^2} \right]_X + \left[ \frac{V\Psi}{C^2} \right]_y + (DA)X + O(\varepsilon^2), \]

\[ \kappa D\Omega CA + iD\Omega y CB - \kappa^2 \Psi = i\varepsilon\kappa D \left[ \frac{\Psi}{D} \right]_X + (UA)_X + VA_y + O(\varepsilon^2), \tag{4.4a} \]

and boundary conditions

\[ i\omega Z_h(B - \varepsilon h X A) = \left( i\Omega C + \varepsilon U \frac{\partial}{\partial X} + \varepsilon V \frac{\partial}{\partial y} - \varepsilon V_y + \varepsilon h X U_y \right) \Psi + O(\varepsilon^2), \]

\[ -i\omega Z_g(B - \varepsilon g X A) = \left( i\Omega C + \varepsilon U \frac{\partial}{\partial X} + \varepsilon V \frac{\partial}{\partial y} - \varepsilon V_y + \varepsilon g X U_y \right) \Psi + O(\varepsilon^2). \tag{4.4b} \]

If we add the first two equations of (4.4a) to the \( y \)-derivative, multiplied by \( \Omega^2 \), of the third equation, we obtain

\[ \Omega^2 \left( \frac{\Psi_y}{\Omega^2} \right)_y + (\Omega^2 - \kappa^2) \Psi = i\varepsilon\Omega^2 \left[ \frac{C}{\Omega} \left( \frac{U\Psi}{C^2} + DA \right)_X + \frac{\kappa D}{\Omega^2} \left( \frac{\Psi}{D} + UA \right)_X \right. \]

\[ + \frac{1}{\Omega C} (V \Psi)_y + \frac{\kappa D V}{\Omega^2} A_y + iD \left( \frac{1}{\Omega^2} \left[ UB_X + (VB)_y \right] \right) \] \[ \left. + O(\varepsilon^2) \right) \tag{4.5} \]

while

\[ A = \frac{\kappa \Psi}{D\Omega C} - \frac{U_y \Psi_y}{D\Omega^2 C^2} + O(\varepsilon) \]

\[ B = i \frac{\Psi_y}{D\Omega C} - \frac{\varepsilon}{\Omega C} \left[ U \left( \frac{\Psi_y}{D\Omega C} \right)_X + \left( \frac{V\Psi_y}{D\Omega C} \right)_y \right] + O(\varepsilon^2). \tag{4.6} \]

We expand to \( O(\varepsilon^2) \)

\[ A = A_0 + \varepsilon A_1 + \ldots, \quad B = B_0 + \varepsilon B_1 + \ldots, \quad \Psi = \Psi_0 + \varepsilon \Psi_1 + \ldots, \quad \kappa = \kappa_0 + \ldots. \]

There is no need for a \( \kappa_1 \)-term. For notational convenience we retain \( \kappa \) in the formulas. To leading order equation (4.5) reduces to the Pridmore-Brown eigenvalue equation

\[ \Omega^2 \left( \frac{\Psi_{0y}}{\Omega^2} \right)_y + (\Omega^2 - \kappa^2) \Psi_0 = 0, \tag{4.7a} \]

with boundary conditions

\[ \Psi_{0y} = \frac{D\Omega^2 C^2}{i\omega Z_h} \Psi_0, \quad \Psi_{0y} = -\frac{D\Omega^2 C^2}{i\omega Z_g} \Psi_0, \tag{4.7b} \]

with \( \Psi_0 \) and \( \kappa \) to be determined. Through \( g, h, U \) and \( C, \Psi_0 \) and \( \kappa \) will vary in \( X \). Since the equation and boundary conditions are linear and homogeneous, \( \Psi_0 \) is only determined up to a slowly varying amplitude. In the following we will derive an adiabatic invariant for \( \Psi_0 \) to determine this slowly varying factor.
4.2. Final solution: the general case

The next order of the expanded equation (4.5) is the inhomogeneous Pridmore-Brown equation in \( \Psi_1 \)

\[
\Omega^2 \left( \frac{\Psi_1_y}{\Omega^2} \right)_y + (\Omega^2 - \kappa^2) \Psi_1 = i \Omega^2 \left[ C \left( (\kappa + U \Omega/C) \frac{\Psi_0}{\Omega^2} - U_y \Psi_{0y} \right)_x + \frac{\kappa^2 V \Psi_{0y}}{\Omega^2 C} \right] + \kappa D \frac{\omega \Psi_0}{D \Omega C} - \frac{U U_y \Psi_{0y}}{D \Omega^2 C^2} + \left( V \frac{\Psi_0}{\Omega C} - \frac{D U}{\Omega^2 C^2} \left( \frac{\Psi_{0y}}{\Omega^2} \right)_x - \frac{1}{\Omega^2 C} \left( \frac{V \Psi_{0y}}{\Omega} \right)_y \right] \quad (4.8a)
\]

with boundary conditions

\[
i \left( \Psi_{1y} - \frac{D \Omega^2 C^2}{i \omega Z_h} \Psi_1 \right) = \frac{DU}{\Psi_0} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_x + \frac{DV}{\Psi_0} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_y + h_x \kappa \Psi_0,
\]

\[
i \left( \Psi_{1y} + \frac{D \Omega^2 C^2}{i \omega Z_g} \Psi_1 \right) = \frac{DU}{\Psi_0} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_x + \frac{DV}{\Psi_0} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_y + g_x \kappa \Psi_0. \quad (4.8b)
\]

We do not aim to solve this equation, as it would involve again undetermined homogeneous solutions. Instead, we derive a solvability condition on \( \Psi_0 \), required for \( \Psi_1 \) to exist. For this we combine equations (4.8a) and (4.7a) as follows

\[
i \int_g^h \frac{1}{D \Omega^2 C^2} \left[ \text{[equation 4.8a] } \Psi_0 - \text{[equation 4.7a] } \Psi_1 \right] \, dy \quad (4.9)
\]

to eliminate \( \Psi_1 \), and apply the boundary conditions (4.8b) to obtain

\[
\left[ \left( \frac{U}{\Omega^2 C^2} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_x + \frac{V}{\Omega^2 C^2} \left( \frac{\Psi_0 \Psi_{0y}}{D \Omega C} \right)_y + \frac{V \kappa \Psi_0^2}{U D \Omega^2 C^2} \right)_g^h \right.
\]

\[
- \int_g^h \frac{\Psi_0}{D \Omega C} \left[ C \left( (\kappa + U \Omega/C) \frac{\Psi_0}{\Omega^2} - U_y \Psi_{0y} \right)_x + \frac{\kappa^2 V \Psi_{0y}}{\Omega^2 C} \right]_g^h
\]

\[
+ \kappa D \frac{\omega \Psi_0}{D \Omega C} - \frac{U U_y \Psi_{0y}}{D \Omega^2 C^2} + \left( \frac{V \Psi_0}{\Omega C} - \frac{D U}{\Omega^2 C^2} \left( \frac{\Psi_{0y}}{\Omega^2} \right)_x - \frac{1}{\Omega^2 C} \left( \frac{V \Psi_{0y}}{\Omega} \right)_y \right)_g^h \quad (4.10a)
\]

By several partial integrations, using the defining differential equation and collecting together the \( X \)-derivative of \((\kappa + U \Omega/C)\Psi_0^2/D \Omega^2 C^2 - \omega U_y \Psi_0 \Psi_{0y}/D \Omega^4 C^4\), we obtain after a considerable amount of algebra the equation

\[
\left[ \left( \frac{\Psi_0}{D \Omega^2 C^2 \Psi_0} \left( \frac{U \Psi_0^2}{\Omega C} \right)_x + \frac{\Psi_{0y}}{D \Omega^2 C^2} \left( \frac{U \Psi_0^2}{\Omega C} \right)_y \right)_g^h \right.
\]

\[
+ \frac{d}{dX} \int_g^h \left( \kappa + \frac{U \Omega}{C} \right) \frac{\Psi_0^2}{D \Omega^2 C^2} - \omega U_y \Psi_0 \Psi_{0y} \frac{U_y}{D \Omega^4 C^4} \, dy = \int_g^h \frac{\kappa U \Psi_0 \Psi_{0y}}{\Omega^4 C^4} \left( \frac{U_y}{D} \right)_x \, dy. \quad (4.10a)
\]

With boundary conditions (4.7b) this becomes

\[
\frac{d}{dX} \left[ \frac{1}{i \omega Z_h} \frac{U \Psi_0^2}{\Omega C} \right]_g^h + \frac{1}{i \omega Z_g} \frac{U \Psi_0^2}{\Omega C} + \int_g^h \left( \kappa + \frac{U \Omega}{C} \right) \frac{\Psi_0^2}{D \Omega^2 C^2} - \omega U_y \Psi_0 \Psi_{0y} \, dy = \int_g^h \frac{\kappa U \Psi_0 \Psi_{0y}}{\Omega^4 C^4} \left( \frac{U_y}{D} \right)_x \, dy \quad (4.10a)
\]
This form is interesting because of its congruence with the acoustic power (see below) in case of cut-on modes in linearly sheared flow and hard walls. An even more elegant form is found by applying relation (A 1) and absorb the boundary terms into the integral, to find

$$\frac{d}{dX} \left[ \int_{g}^{h} \frac{1}{DQ^3C^3} (\omega \kappa \Psi_0^2 + U \Psi_0^2) \, dy \right] = \int_{g}^{h} \frac{\kappa U \Psi_0 \Psi_0 y}{\Omega^4 C^4} \left( \frac{U_y}{D} \right)_X \, dy \quad (4.10b)$$

If we normalise \( \Psi_0 \) to \( \overline{\Psi}_0 \) in some smooth way\( ^\dagger \), and write the slowly varying amplitude \( N \) such that \( \Psi_0(X, y) = N(X) \overline{\Psi}_0(X, y) \), then the equation for \( N \)

$$\frac{d}{dX} [N^2 F] = N^2 G$$

with

$$F(X) = \int_{g}^{h} \frac{1}{DQ^3 C^3} (\omega \kappa \overline{\Psi}_0^2 + U \overline{\Psi}_0^2) \, dy, \quad G(X) = \int_{g}^{h} \frac{\kappa U \overline{\Psi}_0 \overline{\Psi}_0 y}{\Omega^4 C^4} \left( \frac{U_y}{D} \right)_X \, dy$$

leads to a form of an incomplete adiabatic invariant

$$N^2(X) = \frac{N^2(0)}{F(X)} \exp \left( \int_{0}^{X} \frac{G(z)}{F(z)} \, dz \right). \quad (4.13)$$

This is similar to Peake & Cooper (2001); Cooper & Peake (2001); Lloyd & Peake (2013); Rienstra (2016), although here \( F \) and \( G \) are independent of \( X \)-derivatives of \( \overline{\Psi}_0 \).

### 4.3. Final solution: linearly sheared flow

If \( (U_y/D)_X = 0 \), \( G(X) = 0 \) and (4.13) reduces to a complete adiabatic invariant. Since \( D \) is independent of \( y \), this is the case for a linearly sheared mean flow with the property that \( U_y/D = \) a constant. As we found above in (3.9), this is not a restriction. It occurs for any linearly sheared flow that is of slow variation! We have then for a constant \( Q_0 \)

$$\int_{g}^{h} \frac{1}{DQ^3 C^3} (\omega \kappa \Psi_0^2 + U \Psi_0^2) \, dy = Q_0. \quad (4.14)$$

We don’t have to normalise \( \Psi_0 \): the adiabatic invariant is the normalisation. The size of \( Q_0 \) is unimportant in a linear problem, although we can use it to scale the solution conveniently. There is, however, a small reservation. For hard walls we can assume \( \Psi \) to be real, in which case the sign of \( Q_0 \) relates to the propagation direction and cannot be chosen freely. This is due to the fact that for hard-wall real cut-on modes this invariant is, apart from a factor \( \frac{1}{2} \omega \), equal to the acoustic power passing a duct cross section \( X = \) constant, at least, up to the WKB-accuracy. We find for the integral of the time averaged energy flux \( \langle T \rangle \) in \( x \)-direction (2.8b) (with hard walls, \( U_y \) constant in \( y \), real \( \Psi \) and ignoring subscript 0)

$$\mathcal{P}(X) = \int_{g}^{h} \frac{1}{2} \text{Re} \left[ \left( DA + \frac{U \Psi}{C^2} \right) \left( \frac{\Psi}{D} + U A^* \right) \right] \, dy + O(\varepsilon) = \frac{1}{2} \int_{g}^{h} \left( \frac{\kappa \Psi}{DQ C} - \frac{U_y \Psi y}{D^2 C^2} + \frac{U \Psi}{C^2} \right) \left( \frac{\Psi}{D} + \frac{\kappa U \Psi}{DQ C} - \frac{UU_y \Psi y}{D^2 C^2} \right) \, dy + O(\varepsilon) = \frac{1}{2} \int_{g}^{h} \left[ \frac{(\kappa + U \Omega)}{C} \right] \frac{\omega \Psi^2}{DQ^2 C^2} - \frac{\omega^2 U_y \Psi y}{DQ^4 C^4} + \frac{U_y}{2DC^4} \left( \frac{U^2 \Psi y^2}{\Omega^4} \right)_y \, dy + O(\varepsilon) = \frac{1}{2} \omega Q_0 + O(\varepsilon). \quad (4.15)$$

\( ^\dagger \) For example by taking \( F(X) = \pm 1 \).
It is clear that, for $\omega > 0$, $Q_0$ is positive when the mode is right-running, and negative when it is left-running. Another conclusion we may draw is that as far as the left hand side of (4.10) describes the acoustic energy of the mode, it couples with the mean flow, to WKB-accuracy, there where the specific shear $U_y/D$ changes in $X$.

The simplicity of (4.10) is elegant not only physically but also mathematically, because no $X$-derivatives of $\Psi$ and $\Psi_y$ are required. The existence of an adiabatic invariant is a remarkable result, and perhaps not entirely to be expected (Peake & Cooper 2001; Cooper & Peake 2001; Lloyd & Peake 2013; Rienstra 2016), because in general acoustic energy is not conserved in vortical flow. Although for slowly varying real, hard-wall modes in nearly-parallel flow the time-averaged acoustic energy source $\langle D \rangle$ (2.8c) is small†

$$\langle D \rangle = -\frac{i}{2} \text{Re} \left[ i \frac{UU_{yy}}{D^3 C^3} \Psi_y^2 + O(\varepsilon) \right] = 0 + O(\varepsilon),$$

it may be non-zero and its contribution may be finite over a distance $x = O(1/\varepsilon)$.

This completes the leading order solution of the slowly varying modes. From $\Psi_0$ the other amplitudes $A_0$ and $B_0$ follow with (4.6). We end this analysis with checking the consistency of the resulting equation for $\Psi$ with the corresponding results for the potential flow problem, as considered in (Rienstra 1999, 2003b). Without shear, i.e. $U_y = 0$, the mean flow is irrotational and can be given by a potential. $U$ and so $\Omega$ are then functions of $X$ only. If also the sound field is given by a potential $\Phi$ (no vortical perturbations like vorticity shed from sharp edges), such that

$$A = -i\kappa \Phi + \varepsilon \Phi_X, \quad B = \Phi_y, \quad \Psi = -iD\Omega C \Phi - \varepsilon D(U\Phi_X + V\Phi_y) + O(\varepsilon^2)$$

and boundary conditions

$$\Phi_y - \frac{D\Omega^2 C^2}{i\omega Z_h} \Phi = -i\varepsilon h_X \kappa \Phi - \frac{\varepsilon}{\omega Z_h} \left[ U(D\Omega C)^2 X + V(D\Omega C\Phi^2)_y + (DU)_X \Omega C \Phi^2 \right]$$

$$\Phi_y + \frac{D\Omega^2 C^2}{i\omega Z_g} \Phi = -i\varepsilon g_X \kappa \Phi + \frac{\varepsilon}{\omega Z_g} \left[ U(D\Omega C\Phi^2)_X + V(D\Omega C\Phi^2)_y + (DU)_X \Omega C \Phi^2 \right]$$

we find after substitution into (4.5)

$$\Phi_{yy} + (\Omega^2 - \kappa^2) \Phi = \frac{i\varepsilon}{D\Phi} \left[ \frac{\partial}{\partial X} \left( \left( \kappa + \frac{U\Omega}{C} \right) D\Phi^2 \right) + \frac{\partial}{\partial y} \left( \frac{V\Omega}{C} D\Phi^2 \right) \right] + O(\varepsilon^2).$$

After expanding $\Phi = \Phi_0 + \varepsilon \Phi_1 + O(\varepsilon^2)$, this leads in the usual way (cf. Rienstra (1999)) to the adiabatic invariant

$$\frac{d}{dX} \left[ \frac{D^2 U \Omega C}{i\omega Z_h} \Phi_0^2 \right]_h + \frac{D^2 U \Omega C}{i\omega Z_g} \Phi_0^2 \bigg|_g + \int_g^h \left( \kappa + \frac{U\Omega}{C} \right) D\Phi_0^2 \, dy = 0. \quad (4.16)$$

This agrees with (4.10) for $U_y = 0$ if we identify $\Phi_0 = i\Psi_0/D\Omega C$.

† For cut-on hard-wall modes $(v, p, \rho) e^{i\omega t - i\kappa x}$ in parallel shear flow $V = U(y, z) e_x$, this is in general the case. If $p$ is real, then $\rho$ and $u$ are real, while $v$ and $w$ are imaginary. Then $\xi_3$ and $\xi_3$ are real and $D\overline{V} \cdot (\xi \times v^*)$ is imaginary. Since $\Xi = (0, U_z, -U_y)$ is $\Xi \times V = (0, -UU_y, -UU_z)$, and $\rho v^* \cdot (\Xi \times V)$ is imaginary. Eventually, $\langle D \rangle = 0$. 
5. Numerical evaluation of practical examples

5.1. Numerical solution of the Pridmore-Brown equation

The heart of a numerical evaluation of the above solution is the solution of the Pridmore-Brown eigenvalue equation in $y$ (4.7a, 4.7b) along a grid of points in $x$. At each $x$, we consider eigenfunction $\Psi(X, y)$ as a function of $y$ only, and keep $X$ fixed.

For the special case of linear shear flow and uniform soundspeed, it can be solved exactly by Weber’s Parabolic Cylinder functions (Goldstein & Rice 1973). However, this solution, being the difference of exponentially large functions, appeared difficult to evaluate accurately (Rienstra 2016), and is in practice only of use for a limited parameter range. Especially the upstream running modes deteriorate quickly. A numerical solution is therefore inevitable.

A number of methods have been reported in the literature, e.g. Chebyshev spectral collocation methods acting on the vector of linearised Euler equations (Cooper & Peake 2019), which is a classic Galerkin approach based on a symmetric bilinear form.

Not because it’s necessarily faster (we haven’t investigated this) but because of its simple and compact form, we like to use here a method, published recently in (Rienstra 2019), which is a Gauss-Legendre method, which is efficient and needs relatively little points for a high accuracy (unless we have a mean flow with very thin boundary layers).

We rewrite equation (4.7a) with boundary conditions $i\omega \zeta_h \Psi'(h) = \Omega^2 \Psi'(h)$ and $i\omega \zeta_g \Psi'(g) = -\Omega^2 \Psi(g)$ in weak form, which means that we search for eigenfunction $\Psi$ and eigenvalue $\kappa$ such that for every test function $w$

$$\frac{\Psi w}{i\omega \zeta_h} \bigg|_h + \frac{\Psi w}{i\omega \zeta_g} \bigg|_g + \int_g^h -\frac{1}{\Omega^2} \Psi' w' + \left( 1 - \frac{\kappa^2}{\Omega^2} \right) \Psi w \, dy = 0. \quad (5.1)$$

The prime indicates a derivative to $y$. We assume that $\Omega \neq 0$ for any $y$ along the integration interval (a critical layer), to avoid the relatively unimportant critical layer solutions, not studied here; see (Brambley, Rienstra & Darau 2012). $\Psi$ will only be defined up to a multiplicative factor, to be determined below by the adiabatic invariant.

If we assume that $\Psi$ can be written as a sum over a function basis $\{\phi_n\}$ (here: Chebyshev polynomials),

$$\Psi = \sum_{\mu=0}^{\infty} a_\mu \phi_\mu, \quad (5.2)$$

and we use the same basis for the test functions, i.e. $w = \phi_\nu$, $\nu = 0, 1, \ldots$, then (5.1) becomes equivalent with the matrix equation

$$\mathcal{M}(\kappa) \mathbf{a} = 0, \quad \mathbf{a} = (a_0, a_1, \ldots)^T, \quad (5.3)$$

for the (chopped off) doubly infinite symmetric matrix $\mathcal{M}$ with elements

$$\mathcal{M}_{\mu\nu}(\kappa) = \frac{\phi_\mu \phi_\nu}{i\omega \zeta_h} \bigg|_h + \frac{\phi_\mu \phi_\nu}{i\omega \zeta_g} \bigg|_g + \int_g^h \left\{ -\frac{1}{\Omega^2} \phi_\mu \phi_\nu + \left( 1 - \frac{\kappa^2}{\Omega^2} \right) \phi_\mu \phi_\nu \right\} \, dy. \quad (5.4)$$

The $\kappa$-derivative $\mathcal{M}'$ can be expressed in a similar way. The integrals of $\mathcal{M}$ and $\mathcal{M}'$ are evaluated by the Gauss-Legendre method, which is efficient and needs relatively little points for a high accuracy (unless we have a mean flow with very thin boundary layers).

By the following Newton-like iteration in $\kappa$ (Mehrmann & Voss 2004), we search for a solution $\kappa, \mathbf{a}$ that yields a non-trivial null-space for $\mathcal{M}$. Construct the $j+1$-st iterated eigenvalue $\kappa_{j+1}$ with vector $\mathbf{a}_{j+1}$ by solving the approximate equation

$$\mathcal{M}(\kappa_{j+1}) \mathbf{a}_{j+1} \simeq \left\{ \mathcal{M}(\kappa_j) + (\kappa_{j+1} - \kappa_j) \mathcal{M}'(\kappa_j) \right\} \mathbf{a}_{j+1} = 0 \quad (5.5)$$
as follows. Solve the generalised eigenvalue problem

\[ M(\kappa_j) x = \lambda M'(\kappa_j) x, \]

with solutions \( \lambda_1, \lambda_2, \ldots \) and \( x_1, x_2, \ldots \). Choose from all \( \lambda \)'s the absolutely smallest, \( \lambda_{\text{min}} \), and identify \( \kappa_{j+1} = \kappa_j - \lambda_{\text{min}} \) and \( a_{j+1} = x \), where \( x \) corresponds to this \( \lambda_{\text{min}} \). Iterate until \( \| M(\kappa_j) a_j \| \) is small enough. This simple and elegant iteration converges quadratically. Along the slowly varying duct, each previous \( X \)-position provides an excellent starting value, making in general no more than one iteration necessary. Since we deal with one mode, it is no loss of efficiency to find one eigenvalue at a time. 

By using the same Gauss-Legendre points for integral (4.10) as for the elements of \( M \) and \( M' \), we can normalise the found solutions as required. Here we will take the transmitted power at the left entrance of the duct equal to \( \mathcal{P} = 1 \) for right-running modes and at the right entrance equal to \( \mathcal{P} = -1 \) for left-running modes. Note that the normalisation is on \( \Psi^2 \), rather than \( \Psi \), and we have to select a sign for \( \Psi \) at each \( x \)-position. The correct sign is of course such that \( \Psi \) is continuous in \( x \).

The integrals over \( \kappa \) and \( G/F \) may be constructed by incremental trapezoidal integration.

5.2. A practical example

A typical and relevant example, based on a linearly sheared mean flow, is pictured below. The right-running mean flow is described by the problem parameters

\[ \lambda = 0.5, \quad \gamma = 1.4, \quad D_{\text{in}} = 1, \quad \tau_{\text{in}} = 0.2, \quad \mathcal{F} = 0.4496, \quad E = 2.52, \]

and the duct geometry (with inlet at \( x = -3 \)) by

\[ h(x) = 1 - \frac{1}{8}(1 + \tanh(x)), \quad g(x) = 0, \quad -3 < x < 3, \]

giving an estimated \( \varepsilon \approx 0.1 \). See figures 2 for the \( U \) and \( V \) components. The decreasing duct height results into an increase of \( U \) with a slight decrease of \( U_y = \sigma \). \( V \) is obviously at its maximum near \( x = 0 \) at the upper wall.

For the acoustic part we consider a hard wall and a soft wall configuration, both with \( \omega = 15 \). The hard wall case has in total 10 cut-on modes \( p_1, \cdots, p_5 \) right-running and \( p_{-1}, \cdots, p_{-5} \) left-running. The modes \( p_6 \) and \( p_{-6} \) start cut-on at \( x = -3 \) but become cut-off near \( x = -0.055 \). Since we haven’t included a turning point analysis (Rienstra 2000), we excluded these modes. For the soft wall case with \( Z_h = Z_g = 1 + 2i \) we considered the right-running \( p_2 \) and left-running \( p_{-2} \).

With matrix \( M \) chopped off to just \( 21 \times 21 \) (21 Chebyshev terms) and for the integrals using only 40 Gauss-Legendre points we obtain all eigenvalues correct in 11 digits, as can be estimated by doubling these numbers a few times. We expect the same accuracy for the eigenfunctions. The number of \( \kappa \)-iterations, using the result of the previous \( X \)-position as starting value, is typically 1 (depending on the number of \( X \)-positions; here 2400). Per \( X \) and 400 plotpoints in \( y \), this takes altogether about 0.04 seconds (Matlab2018, Intel Core i5-4670 CPU @ 3.4 GHz).

We see that the effects of the sheared mean flow properties are various; see figures (3,4) with snapshots of the pressure distribution in the hard walled duct. For the lowest modal orders, especially upstream running, the asymmetry of the flow is reflected in asymmetric mode shapes. The right-running modes exist only near the lower wall \( y = 0 \) and are negligible at the other side. The upstream running modes behave the same upside down. This behaviour is due to the sign of \( \Omega^2 - \kappa^2 \); the mode is oscillatory in \( y \) when \( \Omega^2 - \kappa^2 \) is positive, and exponentially decaying otherwise. For the present, linear, mean
flow profile the $y$-interval with oscillatory behaviour is either the full interval, or a part adjacent to one of the borders.

In the figures (6,7) the mode shapes in $y$ for various $x$ are given in detail. Indeed modes $p_1, p_{-1}, p_{-2}, p_{-3}$, are practically vanishing at the upper, respectively lower wall. As a result, these modes will not change if the respective wall is lined. Any attenuation due to lining will have to come from the opposite wall.

The modal wave numbers $\kappa_n$ as function of $x$ are given in figure 5. The analogue of the Doppler effect in uniform flow, of increasing wave numbers for upstream running modes and decreasing wave numbers on downstream running modes, is clearly present.

Near $x = 0.4$, the value of $\kappa_5$ changes sign, which is visible in figure 3(e), where the modal wave length suddenly becomes large, and (this is clearer in an animation) the phase velocity changes from positive to negative. The modes seems to propagate to the right for $x < 0.4$ and to the left for $x > 0.4$. This, however, is an optical illusion. The energy of the mode remains right-running everywhere.

As the selected impedance does not have much attenuation, the corresponding figures for the lined wall case are comparable. The snap shots of the pressure distribution, figure 8, show the expected Doppler compression upstream and stretching downstream, and the decay to the right for the right-running mode $p_2$ and to the left for the left-running mode $p_{-2}$. We see here and in figure 9 again that the field of $p_{-2}$ practically vanishes near the lower wall.

The attenuation is better displayed by the transmitted power $10 \log |\mathcal{P}|$ in dB, shown in figure 10. Evidently, the decay is dominated by the (integrated) imaginary parts of $\text{Im} \kappa_{\pm 2}$ (figure 11).

6. Conclusions

The theory of slowly varying modes of WKB type in lined ducts with mean flow is fairly well established for irrotational (potential) mean flow, because the slowly varying mean flow equations can be solved practically analytically and the modal amplitude equations
Figure 3: Pressure field snapshots (order 1 \cdots 5) for $\omega = 15$ and $P = 1$ (right-running)
Figure 4: Pressure field snapshots (order $-1 \cdots -5$) for $\omega = 15$ and $P = -1$ (left-running)
Slowly Varying Modes in Shear Flow

Figure 5: Axial wave numbers $\kappa_1, \ldots, \kappa_{-1}$ (from top to bottom) as function of $x$. $\kappa_6$ and $\kappa_{-6}$ become cut-off and merge together at $x = -0.055$.

Figure 6: Right-running modal functions $\Psi_n(x,y)$ for $\omega = 15$, as function of $y - g$ for varying $x$. $\Psi_6$ becomes cut-off and is not included.
Figure 7: Left-running modal functions $\Psi_n(x,y)$ for $\omega = 15$, as function of $y - g$ for varying $x$. $\Psi_{-6}$ becomes cut-off and is not included.

can be integrated completely to a conserved quantity called adiabatic invariant. This is not the case for a rotational mean flow, at least not in 3D cylindrical ducts. The equations for the mean flow as well as the modal amplitudes require a numerical integration, and no simple adiabatic invariant is available.

We found here that the simplifications due to a 2D duct provide conditions that allow more progress analytically. For a general sheared mean flow, the mean flow and amplitude equations cannot be solved without a numerical integration but the amplitudes satisfy an equation that is almost an adiabatic invariant. If the specific mean vorticity remains constant in axial direction (rather than constant along streamlines, what it does in general), the *incomplete* adiabatic invariant reduces to a *complete* one.

This is just what happens for the particular case of a linearly sheared mean flow. Even more, not only the specific mean vorticity is constant, but the whole flow field can be determined practically analytically (up to an algebraic equation).

In conclusion, there is a remarkable, but apparently uncorrelated coincidence for 2D modes in slowly varying linearly sheared flow: (1) the prevailing Pridmore-Brown equation can be solved analytically; (2) the mean flow equations can be solved (practically) analytically; (3) the modal amplitude equation can be integrated completely.

Unfortunately, the analytical solution of the Pridmore-Brown equation, being the
Slowly Varying Modes in Shear Flow

\[ \kappa_2 = 9.56 - 0.84i \ldots 7.99 - 0.13i \]

\[ \kappa_2 = -30.25 + 0.38i \ldots - 61.22 + 0.53i \]

Figure 8: Pressure field snapshots (orders ±2) for \( \omega = 15 \) and \( Z = 1 + 2i \)

difference of exponentially large functions, did not prove to be very useful here. The simplicity of the other two results, however, make the linear shear flow problem comparable to the potential flow problem.

Appendix A

A useful relation for sufficiently smooth function \( G \) and solutions \( \Psi, \kappa \) of the Pridmore-Brown eigenvalue equation (4.7a) is the integral

\[ \int_g^h \Psi^2 G \, dy = \int_g^h \left[ \frac{G}{\Omega^2} (\kappa^2 \Psi^2 + \Psi_y^2) + \frac{G_y}{\Omega^2} \Psi \Psi_y \right] \, dy - \frac{G \Psi \Psi_y}{\Omega^2} \Big|_g^h. \]  

(A 1)

We prove it by noting that from the defining equation we have

\[ \Psi^2 G = \frac{\kappa^2}{\Omega^2} \Psi^2 G - \left( \frac{\Psi_y}{\Omega^2} \right)_y \Psi G \]

and use partial integration. In particular, with \( G = U/(D \Omega C^3) \), \( G_y = \omega U_y/(D \Omega^2 C^4) \)

\[ \int_g^h \left( \kappa + \frac{U \Omega}{C} \right) \frac{\Psi^2}{D \Omega^2 C^2} \, dy - \frac{\omega U_y}{D \Omega^4 C^4} \Psi \Psi_y \, dy = \]

\[ \int_g^h \left( \frac{\kappa}{D \Omega^2 C^2} \Psi^2 \, dy + \int_g^h \frac{U}{D \Omega^3 C^3} \Psi^2 - \frac{\omega U_y}{D \Omega^4 C^4} \Psi \Psi_y \, dy = \right. \]

\[ \left. \int_g^h \frac{\kappa}{D \Omega^2 C^2} \Psi^2 \, dy + \int_g^h \frac{U}{D \Omega^3 C^3} (\kappa^2 \Psi^2 + \Psi_y^2) \, dy - \left[ \frac{U}{D \Omega^3 C^3} \Psi \Psi_y \right]_g^h = \right. \]

\[ \left. \int_g^h \frac{1}{D \Omega^3 C^3} (\omega \kappa \Psi^2 + U \Psi_y^2) \, dy - \left[ \frac{U}{D \Omega^3 C^3} \Psi \Psi_y \right]_g^h \right] \]
Figure 9: Real and imaginary parts of modal functions $\Psi_{\pm 2}(x, y)$ for $\omega = 15$, $Z = 1 + 2i$, as function of $y - g$ for varying $x$.

Figure 10: Axial acoustic power $10 \log |P|$ in dB of right-running mode 2 and left-running mode $-2$ as function of $x$. 
Figure 11: Complex axial wave numbers $\kappa_{\pm 2}$ varying from $x = -3$ (■) to $x = 3$ (○).

Declaration of Interests

The author reports no conflict of interest.

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