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COMPUTATION OF LAMINAR VISCOUS FLOWS USING VON MISES COORDINATES

by

Paul Stephen Carson

A Dissertation
Submitted to the Faculty of Graduate Studies and Research
Through the Department of Mathematics & Statistics
in Partial Fulfillment
of the Requirements for the degree of
Doctor of Philosophy
at the University of Windsor

Windsor, Ontario, Canada

1994
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ABSTRACT

In this thesis, we address the problem of computing laminar incompressible viscous flows. For such flows, there are various possibilities for the formulation of the problem. These include primitive variable (velocity and pressure), velocity-vorticity and stream function-vorticity formulations. Each has its own strengths and weaknesses.

In the stream function-vorticity formulation for two-dimensional flow, the difficulty is primarily associated with determination of vorticity at a boundary. In this thesis, we employ a variation of the stream function-vorticity formulation whereby incompressible, viscous flow in an internal complex geometry is formulated in terms of von Mises coordinates. That is, stream function is used as an independent rather than dependent variable. This formulation provides a rectangular computational domain with both Dirichlet and von Neumann boundary conditions for unknown functions, the vertical cartesian coordinate and the vorticity, in terms of the horizontal cartesian coordinate and the stream function. The governing second order nonlinear partial differential equations are solved by SLOR on uniform and, if required, clustered grids. A number of procedures for surmounting the problem of determining vorticity at a boundary are available. A novel approach to this problem is applied in this thesis.

A difficult, but well-documented, test problem was chosen to study the applicability of the von Mises formulation in viscous flows. It has been shown that extreme care must be taken in applying von Mises coordinates to viscous flow situations. In particular, viscosity is known to generate vorticity in the flow field and to cause, under appropriate conditions, flow separation. Of these two phenomenon, rotational flow and viscous separation, it is shown that rotational effects can be handled with no more
difficulty than experienced by conventional methods. However, separation cannot be handled directly by the von Mises formulation, and erroneous results may be obtained if not used carefully. This presents a challenging problem which has been overcome by developing an innovative way to predict the location of the streamline which divides the main flow from the recirculating region. In this way, the von Mises formulation can be used to study separated 2D viscous flows. This approach is used to predict the re-attachment length for the flow over a backward facing step and the results, when compared to other numerical data, confirm the applicability and accuracy of the method.
DEDICATION

To the memory of my Dad. I don’t think he knew what he was letting himself in for when he told me to stay in school. I also don’t think he meant for me to do it for as long as I have. Many of my friends have told me they wish they had parents like mine, not perfect, but there when you needed them. And I needed them! Thanks.
ACKNOWLEDGEMENTS

I owe everyone. They know who they are, so there is no point in trying to list them all. Maybe one day, if I keep paying back what I have received, I will go out even. Only God keeps that score.
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LIST OF SYMBOLS OR NOMENCLATURE

E, F, G  coefficients of First Fundamental Form
h  energy function
J  Jacobian
L  characteristic length
p  pressure function
q  speed function
r  radial component of polar coordinate
U_∞  constant flow speed at infinity
u  velocity component in x direction
v  velocity component in y direction
W  ±J
x  cartesian coordinate
x_{LE}  value of x at the leading edge
x_{TB}  value of x at the trailing edge
y  cartesian coordinate

Greek letters

α  local angle of inclination of streamline with x-axis
ε  error
η  stretching coordinate
θ  angular component of polar coordinate, angle between coordinate curves in (φ,ψ) net
κ  curvature
\( \mu \) constant viscosity

\( \xi \) stretching coordinate

\( \rho \) constant fluid density

\( \phi \) curvilinear coordinate

\( \psi \) curvilinear coordinate, stream function

\( \omega \) vorticity function
CHAPTER I
INTRODUCTION

The basic tools used by engineering and mathematical researchers to gain an understanding of the various physical phenomena associated with the dynamics of fluid flow are mathematical analysis (exact or non-exact solution methods), experimental (testing) and computational methods. Before the development of high-speed computers, experimentalists developed various simple devices to make measurements of flow quantities of particular interest to them. Analytical researchers then tried to duplicate these experimental results using fairly simple mathematical models and analysis to provide further guidelines to the experimentalists. Then the process was repeated, sometimes with the analysis leading the experimentation.

As the design of modern high-speed computers has become more and more sophisticated, there has been a greater demand for more detailed and accurate numerical analysis of the flow fields under investigation instead of using costly experimental methods, e.g., wind tunnel testing and prototype development. This need has made computational methods more than an equal partner with experimental methods in efforts to analyze complex fluid flow geometries.

Prior to about 1970, many papers appeared on the numerical solution of the Navier-Stokes equations governing the flow of a viscous incompressible fluid in an internal complex geometry. See, for example, a list of over 300 papers in the Ph.D. dissertation by P.J. Roache [I.1]. In the past ten years or so, more authors have taken an interest in this type of problem. For example, in order to stimulate a fruitful debate among computational fluid dynamics (CFD) specialists and to assess the capabilities of
various numerical methods to deal with laminar flows in complex geometries, the International Association of Hydraulic Research (IAHR) Working Group on Refined Modelling of Flows decided to devote its Fifth Meeting to this subject in 1982. This meeting will be discussed in detail in Chapter III of this thesis.

There are various possibilities for the formulation of the problem of viscous incompressible flow. These include primitive variable (velocity and pressure), velocity-vorticity and stream function-vorticity formulations. The general solution procedure consists of discretizing the differential equations and boundary conditions over the fluid flow region and solving the resulting system of algebraic equations. Finite difference methods are employed in the discretization in this thesis. In general, when the full Navier-Stokes equations in terms of velocity and pressure (primitive variables) are solved, a specified velocity or velocity gradient (usually zero) must be given at the last downstream station due to the elliptic effect of the streamwise diffusion term. The upstream boundary conditions consist of specified velocity profiles at the upstream location. Along a non-porous body surface, the no-slip boundary conditions \( u = v = 0 \) are applied. For a line of symmetry, the normal velocity component and the normal derivative of the longitudinal velocity component must be zero. A downstream pressure boundary condition is needed. This boundary condition can either be a specified value of the pressure or a specified value of the pressure gradient in the direction of flow. This specification of boundary conditions on pressure can lead to serious numerical difficulties. However, the primitive variable approach offers the fewest complications in extending the calculations to three dimensions.

The velocity-vorticity approach requires the vorticity equation, the continuity equation, and the equations that define vorticity in terms of velocity gradients. A
combination of the continuity equation and the definition of vorticity can be made after differentiation of both equations. This yield elliptic equations for the velocity components. An inconvenience of the velocity-vorticity formulation is that pressure is not directly obtained and consequently additional calculations are required for its determination.

For two-dimensional and for axi-symmetric flows, in order to overcome some of the above difficulties, it is convenient to introduce stream function and vorticity as dependent variables. The equation of continuity is automatically satisfied and the resulting system consists of two coupled nonlinear equations which are solved numerically by some iterative procedure. In the stream function-vorticity formulation, the difficulty is primarily associated with determination of vorticity at a boundary. In this thesis, we employ a variation of the stream function-vorticity formulation whereby the flow is formulated in terms of von Mises coordinates.

Aside from the fact that these coupled equations are nonlinear, there are several other difficulties associated with their solution. One of the major difficulties is that the values of vorticity on no-slip boundaries are not known a priori, while these values are needed in order to solve the discretized problem. In terms of von Mises coordinates another major difficulty is specification of the vertical cartesian coordinate, once the transformation has been made, on dividing streamlines between recirculating and non-recirculating flow and on free surfaces. The flow domain of the problem can also introduce other additional difficulties in the numerical method. Some authors prefer the use of velocity-pressure formulation of the Navier-Stokes equations in order to avoid the difficulties arising from the introduction of vorticity. However, as noted earlier, the pressure equation is complicated and introduces its own additional difficulties.
In order to test a numerical method, it is customary to choose a simple model problem. Whatever choice is made mathematically, it is usually non-trivial because of singularities that often arise at corners. Several numerical methods have been proposed in the literature differing in the choice of discretization schemes, the boundary approximations used to define vorticity on no-slip walls and the methods used to solve the resulting system of algebraic equations. Whatever test problem is chosen to be solved by whatever computational method, the numerical solutions are usually compared in terms of the values of stream function or vorticity at some representative points and also by comparing the values of certain parameters of the flow. Both techniques are employed in this thesis.

A primary difficulty in most cases is to choose a coordinate frame that simplifies both the correlation of measured data and the construction of predictive models. In analyzing fluid motions theoretically, cartesian coordinates are commonly adopted. However, Cartesian coordinates are not always the best choice. For example, the rectangular cartesian coordinate system may not be the appropriate choice since the physical interpretation of many quantities becomes elusive when flow direction and coordinate direction do not coincide. To estimate changes in properties of the flow between two points on a streamline requires integration along a curve, generally a complicated operation. Also, there are practical difficulties in aligning instruments accurately with some externally imposed rectangular frame.

The obvious choice has been to choose curvilinear coordinates. If one coordinate direction can be chosen almost parallel to the mean flow direction, then extra terms arising from deviation of the mean flow from the coordinates may be small enough to be approximated in calculation schemes or ignored in interpretation of measurements.
For instance, it is better to use cylindrical coordinates for flow around a circular cylinder or the flow in a circular pipe and to use spherical coordinates for the flow around a sphere. This is because proper coordinates can be chosen corresponding to the shape of boundaries.

In the early 1970's, Martin [I.2] introduced a natural curvilinear coordinate system $(\phi, \psi)$ in the physical plane $(x, y)$ where $\psi$ is the stream function, to study the geometry of certain steady two-dimensional, incompressible, viscous flows. This formulation has been used by Grossman and Barron [I.3] to numerically investigate incompressible, irrotational, inviscid flow over symmetric airfoils at zero angle of incidence. They have chosen the coordinate system to be orthogonal which, in their case, implies that the curves $\phi(x, y) = \text{constant}$ are potential curves. They found that it was not possible to determine analytically where the leading and trailing edges are mapped into the $(\phi, \psi)$ system and that numerically obtained values for the leading and trailing edges are not very accurate, resulting in inaccuracies in the solution near these points. During a further study of incompressible, irrotational, inviscid flows, Barron [I.4] introduced von Mises coordinates $(x, \psi)$. Using these independent variables, one knows exactly where the leading and trailing edges are mapped in the $(x, \psi)$ plane and inaccuracies in the solution near these points can be eliminated [I.5].

The purpose of this dissertation is to study the feasibility and advisability of applying stream function coordinate methods to viscous flow problems. This dissertation extends Barron's approach from that of two-dimensional, steady, incompressible, irrotational, inviscid flows to the case of viscous flows. In Barron's approach the coordinate $\psi$ is taken as the stream function for the flow being considered. This approach automatically provides a rectangular computational domain $(x, \psi)$, and the need
to do grid generation is avoided except possibly in regions of recirculating flow. This is very significant since it reduces the equations to be solved by two. The flow equations are transformed into von Mises coordinates and solved subject to appropriate boundary conditions, which can be formulated as Dirichlet and von Neumann conditions. Since a boundary coincides with a streamline, as long as separation does not occur, the theoretical treatment becomes easier when choosing streamlines as coordinate axes. However, the basic equations become more complicated, which makes them more difficult to solve. This means the method has merits for some types of flow.

Each of the chapters of the dissertation is now briefly described. In Chapter II, the two-dimensional, steady, incompressible, viscous flow equations, first in $(\phi, \psi)$ system and then in the $(x, \psi)$ system (unstretched), are derived. The flow equations are also given in the $(x, \psi)$ system in stretched coordinates $(\xi, \eta)$. A description of the numerical discretization and basic solution algorithm are also provided in Chapter II. Chapter III discusses a well defined test problem, the laminar flow through a non-trivial configuration, namely flow through a smooth expansion channel. Difficulties with flow in the weakly separated region (i.e., small region of recirculation) are discussed. This problem is not fully resolved by the method of solution proposed in Chapter II. Chapter IV presents a solution to a problem in the inviscid limit of zero viscosity. Since the problem of recirculation is not fully resolved with the test problem in Chapter III, this problem is chosen since the flow is still rotational, but no viscosity is allowed. Hence, recirculating flow is eliminated. The problem is hyperbolic-cosine shear flow about a circular cylinder. Chapter V presents the solution to the flow over a backward facing step along with appropriate boundary conditions. The reattachment point of the primary recirculating region is predicted and the calculated results compared to values obtained
experimentally or by other numerical methods. Difficulties with flow near regions of discontinuity resulting in singularities in the flow parameters are discussed.
CHAPTER II
GOVERNING EQUATIONS

2.1 DIFFERENTIAL GEOMETRY

We first refer to some results from differential geometry which are essential for our purposes (cf. [1.2] and [1.4]).

Consider a coordinate transformation between cartesian coordinates \((x,y)\) and some arbitrary curvilinear coordinates \((\phi, \psi)\) defined by

\[
\overline{r} = \overline{r}(\phi, \psi)
\]

(2.1.1)

where \(\overline{r} = (x,y)\), i.e., \(x = x(\phi, \psi)\) \(y = y(\phi, \psi)\) (See Diagram 2.1.1). Assume that the Jacobian \(J\), defined by \(J = |\overline{r}_\phi \times \overline{r}_\psi| = x_\phi y_\psi - x_\psi y_\phi\) is non-zero in the region of interest. The squared element of arc length along any curve can be represented by

\[
ds^2 = d\overline{r} \cdot d\overline{r} = E(\phi, \psi)d\phi^2 + 2F(\phi, \psi)d\phi d\psi + G(\phi, \psi)d\psi^2
\]

(2.1.2)

where

\[
E = \overline{r}_\phi \cdot \overline{r}_\phi - x_\phi^2 + y_\phi^2
\]

\[
F = \overline{r}_\phi \cdot \overline{r}_\psi - x_\phi y_\psi + y_\phi x_\psi
\]

(2.1.3)

\[
G = \overline{r}_\psi \cdot \overline{r}_\psi - x_\psi^2 + y_\psi^2
\]

are the metrics of the space under consideration (coefficients of the first fundamental form).
2.2 \textbf{FLOW EQUATIONS FOR TWO-DIMENSIONAL, STEADY, INCOMPRESSIBLE, VISCOUS FLOW}

The Navier-Stokes and mass conservation flow equations for the two-dimensional, steady (meaning stationary), incompressible (constant density), viscous flow in terms of physical or rectangular coordinates \((x,y)\) are, in dimensional form,

\[
\bar{u}_x + \bar{v}_y = 0 \quad \text{(continuity)} \tag{2.2.1a}
\]

\[
\bar{\rho}(\bar{u}\bar{u}_x + \bar{v}\bar{u}_y) + \bar{p}_x = \mu(\bar{u}_{xx} + \bar{u}_{yy}) \quad \text{(momentum equations)} \tag{2.2.2a}
\]

\[
\bar{\rho}(\bar{u}\bar{v}_x + \bar{v}\bar{v}_y) + \bar{p}_y = \mu(\bar{v}_{xx} + \bar{v}_{yy})
\]

where \(\bar{u}, \bar{v}\) are velocity components in the \(x\) and \(y\) coordinate directions respectively, \(\bar{p}\) is pressure, \(\bar{\rho}\) is the constant density and \(\mu\) is the viscosity. The bar over a variable indicates its dimensional form. Defining a vorticity function \(\bar{\omega} = \bar{\omega}(\bar{x},\bar{y})\) and energy function \(\bar{h} = \bar{h}(\bar{x},\bar{y})\) by

\[
\bar{\omega} = \bar{v}_x - \bar{u}_y \quad \text{(vorticity)} \tag{2.2.3a}
\]

\[
\bar{h} = \frac{1}{2}\bar{\rho}(\bar{u}^2 + \bar{v}^2) + \bar{p} \quad \text{(energy)} \tag{2.2.4a}
\]

equations (2.2.2a) can be written, eliminating pressure \(\bar{p}\), as

\[
\bar{h}_x - \bar{\rho}\bar{v}\bar{\omega} = -\mu \bar{\omega}_y \quad \text{(momentum)} \tag{2.2.2b}
\]

\[
\bar{h}_y + \bar{\rho}\bar{u}\bar{\omega} = \mu\bar{\omega}_x
\]

Equations (2.2.1a), (2.2.3a) and (2.2.2b) constitute a system of four non-linear partial differential equations in four unknown functions: \(\bar{u}\) and \(\bar{v}\) are the velocity components,
\( \vec{\omega} \) is the vorticity and \( \vec{h} \) is the energy. The state equation (\( \vec{\rho} = \text{constant} \)) and energy equation (2.2.4a) along with the preceding system of four equations constitute a complete system of equations.

Nondimensionalizing with respect to a characteristic length \( L \) and speed \( U_\infty \) according to

\[
\begin{align*}
\tilde{x} &= Lx \\
\tilde{y} &= Ly \\
\tilde{u} &= U_\infty u \\
\tilde{v} &= U_\infty v \\
\tilde{h} &= \tilde{\rho} U_\infty^2 h \\
\tilde{\omega} &= U_\infty \omega/L \\
\tilde{p} &= \tilde{\rho} U_\infty^2 p
\end{align*}
\] (2.2.5)

flow equations (2.2.1a), (2.2.3a), (2.2.4a) and (2.2.2b) become

\[
\begin{align*}
&u_x + v_y = 0 \quad \text{(continuity)} \quad (2.2.1b) \\
&h_x - v\omega = -\frac{1}{R_e} \omega_y \quad \text{(momentum)} \quad (2.2.2c) \\
&h_y + u\omega = \frac{1}{R_e} \omega_x \\
&\omega = v_x - u_y \quad \text{(vorticity)} \quad (2.2.3b)
\end{align*}
\]

where,

\[
\begin{align*}
h &= \frac{1}{2}(u^2 + v^2) + p \quad \text{(energy)} \quad (2.2.4b)
\end{align*}
\]
and $Re = \frac{\bar{\rho}U_aL}{\mu}$ is the Reynolds number.

The equation of continuity (2.2.1b) implies the existence of a stream function $\psi(x,y)$ such that

$$
\begin{align*}
    u &= \psi_y \\
    v &= -\psi_x
\end{align*}
$$

(2.2.6)

Appendix A contains a presentation of equations (2.2.2c) and (2.2.3b) using (2.2.6) in an alternate form more familiar in the literature. We will use one of these forms later on in this thesis (equation (A.1)).

Following Martin [1.2], we now proceed to write equations (2.2.1b), (2.2.2c), (2.2.3b) and (2.2.4b) in the $(\phi, \psi)$ curvilinear coordinate system where $\phi = \phi(x,y)$, $\psi = \psi(x,y)$. The essential feature is that rather than consider an arbitrary $(\phi, \psi)$ net, Martin chose $\psi = \text{constant}$ curves to correspond to the streamlines and the function $\psi(x,y)$ to be the stream function defined in equation (2.2.6). Since the flow moves along the streamlines, this is a natural choice of coordinate system on which to do a numerical calculation [1.4]. For the present, the curves $\phi(x,y) = \text{constant}$ are left arbitrary, to be chosen in a convenient manner later. We assume the Jacobian $J \neq 0$ anywhere in the flow region and that the fluid flows along streamlines $\psi = \text{constant}$ in the direction of increasing $\phi$ so that $J > 0$. In this manner the curvilinear coordinate system is more definite, being tied analytically to the actual flow problem. According to Barron [1.4], a second important aspect of Martin's method (to be seen shortly) is that the physical variables $(u, v)$ are replaced by the geometric variables $E$, $F$ and $G$ in the flow equations. Hence, the metric coefficients $E$, $F$ and $G$ are determined as part of the solution of the
flow equations.

Considering E, F, G, h and ω as functions of φ and ψ, flow equations (2.2.1b), (2.2.2c) and (2.2.3b) along with (2.2.4b) are transformed to the (ϕ, ψ) coordinate system.

Martin [1.2] has shown that the continuity equation (2.2.1b) is equivalent to (nondimensionalizing)

\[ q^2 - u^2 + v^2 = \frac{E}{J^2} \]  

(continuity) \hspace{1cm} (2.2.1c)

where \( J^2 = EG-F^2 \).

Hence the energy equation (2.2.4b) can be written as

\[ h = \frac{E}{2J^2} + p \]  

(energy) \hspace{1cm} (2.2.4c)

Martin has also shown that the momentum and vorticity equations (2.2.2c) and (2.2.3b) become (again nondimensionalizing)

\[ Gh_\phi - F(h_\psi + \omega) = - \frac{1}{Re} J\omega_\psi \]  

(momentum) \hspace{1cm} (2.2.2d)

\[ -Fh_\psi + E(h_\psi + \omega) = - \frac{1}{Re} J\omega_\phi \]  

(vorticity) \hspace{1cm} (2.2.3c)

\[ \omega = \frac{1}{J} \left[ \left( \frac{F}{J} \right)_\phi - \left( \frac{E}{J} \right)_\psi \right] \]
Equations (2.2.2d) can be written as (Barron [1.4]) (see Appendix B1),

\[ h_\psi = \frac{1}{R_c} \frac{F}{J} \omega_\phi - \frac{1}{R_c} \frac{E}{J} \omega_\psi \]  
(momentum)  
(2.2.2e)

\[ h_\psi = -\omega + \frac{1}{R_c} \frac{G}{J} \omega_\phi - \frac{1}{R_c} \frac{F}{J} \omega_\psi \]

Finally an equation referred to by Martin as the Gauss equation, states that the Gaussian curvature \( K \) is zero. That is,

\[ K = \frac{1}{J} \left[ \left( \frac{J}{E} \Gamma_{11}^2 \right)_\psi - \left( \frac{J}{E} \Gamma_{22}^2 \right)_\phi \right] = 0 \]  
(Gauss)  
(2.2.7)

where \( \Gamma_{11} \) and \( \Gamma_{12} \) are Christoffel symbols, and

\[ \Gamma_{11}^2 = \frac{-FE_\phi + 2EF_\phi - EE_\psi}{2J^2} \]

\[ \Gamma_{12}^2 = \frac{EG_\psi - FE_\phi}{2J^2} \]

For a plane provided with a curvilinear coordinate system \((\phi, \psi)\), the Gaussian curvature always equals zero.

Equations (2.2.2e), (2.2.3c) and (2.2.7) are four partial differential equations for the five unknown functions \( E, F, G, h \) and \( \omega \). This system is underdetermined because of the arbitrariness of the curves \( \phi(x,y) = \) constant chosen to define the coordinate system. In the next section a choice for \( \phi = \phi(x,y) \) is proposed which provides

---

\(^1\)The reason for including what appear to be simple and obvious derivations in the appendices is either for completeness or because this is the first time the equation appears in the literature.
boundary conditions in a simple form for numerical calculations and removes this arbitrariness. It will be shown that Gauss' equation is automatically satisfied for this choice of $\phi$ since $E$, $F$ and $G$ will then be known in terms of a single unknown function. This will be shown to reduce the four equations eventually to two equations in terms of only two unknown functions, this new function and the vorticity $\omega$.

Of course, other choices are possible (and perhaps desirable in certain flow configurations), such as requiring the grid system to be orthogonal. However, these lead to additional equations to be solved and the boundary conditions are more complicated.

### 2.3 THE VON MISES TRANSFORMATION

As indicated in the previous section, in order to remove the arbitrariness in $\phi = \phi(x,y)$, it is convenient to choose

$$\phi(x,y) = x$$  (2.3.1)

so that the equations of motion (2.2.2e) and (2.2.3e) and the Gauss' equation (2.2.7) are formulated with $(x,\psi)$ as independent variables rather than $(\phi,\psi)$, i.e., von Mises coordinates.

Using equation (2.3.1) in the expressions for the metric coefficients $E$, $F$ and $G$, equations (2.1.3) give the metrics in terms of a single unknown function $y = y(x,\psi)$:

$$E = 1 + y_x^2$$

$$F = y_x y_\psi$$  (2.3.2)

$$G = y_\psi^2$$

The Jacobian $J$ becomes

$$J = y_\psi$$  (2.3.3)
The continuity equation (2.2.1c) becomes, using (2.3.2) and (2.3.3),

\[ q^2 - u^2 + v^2 = \frac{1+y_x^2}{y_\psi} \]  \hspace{1cm} (continuity) \hspace{1cm} (2.3.4)

From (2.2.6),

\[ u = \frac{1}{y_\psi} = \psi_y \]  \hspace{1cm} (2.3.5)

\[ v = \frac{y_x}{y_\psi} - uy_x = -\psi_x \]

Using (2.3.2) and (2.3.3), the energy equation (2.2.4c) becomes

\[ h = \frac{1+y_x^2}{2y_\psi^2} + p \]  \hspace{1cm} (energy) \hspace{1cm} (2.3.6)

The momentum equations (2.2.2e) become, using (2.3.2) and (2.3.3),

\[ h_x = \frac{1}{R_e} \left[ y_x \omega_x - \frac{(1+y_x^2)}{y_\psi} \omega_\psi \right] \]  \hspace{1cm} (momentum) \hspace{1cm} (2.3.7a)

\[ h_\psi = -\omega + \frac{1}{R_e} [y_x \omega_x - y_\psi \omega_\psi] \]

The vorticity equation (2.2.3c) becomes, using (2.3.2) and (2.3.3)

\[ \omega = \frac{1}{y_\psi} \left[ y_{xx} - \left( \frac{1+y_x^2}{y_\psi} \right)_\psi \right] \]  \hspace{1cm} (vorticity) \hspace{1cm} (2.3.8a)

Finally, Gauss' equation (2.2.7) is automatically satisfied since it is equivalent to \( y_{x\psi} = y_{\psi x} \) for all \((x,\psi)\), i.e., the transformation identically satisfies (2.2.7). Since the coordinates \((\phi,\psi)\) satisfy Gauss' equation in order to form a curvilinear net, the von
Mises transformation also must satisfy it so as to form a curvilinear net \((x, \psi)\) (cf. Appendix C).

It should be noted at this point that we have transformed the equations of motion in terms of rectangular cartesian coordinates \((x, y)\) to a curvilinear coordinate system \((\phi, \psi)\) using Martin's method. Then, using the von Mises transformation, we further transform the flow equations to another curvilinear coordinate system \((x, \psi)\).

\[
\begin{align*}
(x, y) & \quad \rightarrow \quad (\phi, \psi) & \quad \rightarrow \quad (x, \psi) \\
\text{physical plane} & & \text{curvilinear plane} & & \text{computational plane}
\end{align*}
\]

This is actually the starting point for this thesis. The streamlines \(\psi = \text{constant}\), which are curved in the physical plane, are mapped to horizontal straight lines in the computational plane.

Equations (2.3.7a) and (2.3.8a) will now be written in a more convenient form to serve as the starting point for the numerical work which will follow.

Expand equation (2.3.8a) to get

\[
y_x^2 \omega = y_x^2 y_{xx} - 2 y_x y_y y_{xy} + \left( 1 + y_x^2 \right) y_{\psi \psi}
\]  

(2.3.8b)

Define the operator

\[L\{ \} = y_x^2 \frac{\partial^2}{\partial x^2} - 2 y_x y_y \frac{\partial^2}{\partial x \partial \psi} + \left( 1 + y_x^2 \right) \frac{\partial^2}{\partial \psi^2}\]

Then, equation (2.3.8b) can be written as

\[L\{y\} - y_x^2 \omega = 0\]
For use later, the above equation is written as

\[ L\{y\} - y_\psi^2 \omega_\psi = 0 \quad \text{(vorticity)} \quad (2.3.8c) \]

where \( y = y(x, \psi) \).

Eliminating \( h \) from equations (2.3.7a), using \( h_{x\psi} = h_{\psi\psi} \), yields

\[ L\{\omega\} - R_x y_\psi \omega_x - y_\psi^2 \omega_\psi = 0 \quad (2.3.7b) \]

where \( \omega = \omega(x, \psi) \) (cf. Appendix D for a derivation of this equation).

Equations (2.3.8c) and (2.3.7b) are two elliptic partial differential equations which must be solved for the two unknown functions \( y = y(x, \psi) \) and \( \omega = \omega(x, \psi) \). The boundary conditions associated with these equations are problem specific and will be discussed in later chapters.

To find the pressure, we use the equations (2.3.7a) to get an equation for energy \( h = h(x, \psi) \) and then use equation (2.3.6) to solve for the pressure \( p = p(x, \psi) \) (cf. Appendix E).  

2.4 **NUMERICAL ALGORITHM**

2.4.1 **Finite Difference Formulation**

The equations will be solved by approximating derivatives by finite differences. Hence, the equations to be solved, i.e., equations (2.3.8c) and (2.3.7b), in difference operator notation, are respectively:

\[
\left[ A_0^{\omega} \delta_{xx} + B_0^{\omega} \delta_{x\psi} + C_0^{\omega} \delta_{\psi\psi} + E_0^{\omega} \delta_\psi \right] y_0^{(n+1)} = 0
\]
and

\[
\left[ A_{ij}^{(n+1)} \delta_{xx} + B_{ij}^{(n+1)} \delta_{x\psi} + C_{ij}^{(n+1)} \delta_{\psi\psi} + \text{Re} D_{ij}^{(n+1)} \delta_x + E_{ij}^{(n+1)} \delta_\psi \right] \omega_{ij}^{(n+1)} = 0
\]

or rewriting these two equations in a more compact notation

\[
\left[ A_{ij}^{(k)} \delta_{xx} + B_{ij}^{(k)} \delta_{x\psi} + C_{ij}^{(k)} \delta_{\psi\psi} + \alpha \text{Re} D_{ij}^{(k)} \delta_x + E_{ij}^{(k)} \delta_\psi \right] \phi_{ij}^{(n+1)} = 0 \tag{2.4.1.1}
\]

where,

\[
\phi = \begin{cases} 
  y(x,\psi) & \text{if } \alpha = 0 \\
  \omega(x,\psi) & \text{if } \alpha = 1 
\end{cases}
\]

are the unknowns and,

\[
k = \begin{cases} 
  n & \text{if } \phi = y \\
  n+1 & \text{if } \phi = \omega \quad (\text{except in } E_{ij}^{(k)} \text{ where } \omega_{ij} \text{ is at (n)})
\end{cases}
\]

is the iteration number.

The operators \( \delta_{xx}, \delta_{x\psi}, \delta_{\psi\psi}, \delta_x \) and \( \delta_\psi \) represent 3-point central difference operators, and are difference approximations to the partial derivatives in (2.3.8c) and (2.3.7b).

The coefficients \( A_{ij}^{(k)}, B_{ij}^{(k)}, C_{ij}^{(k)}, D_{ij}^{(k)} \) and \( E_{ij}^{(k)} \) in (2.4.1.1) are:

\[
\begin{align*}
A_{ij}^{(k)} &= (\delta_x y)_{ij}^2 \\
B_{ij}^{(k)} &= -2(\delta_x y)_{ij}(\delta_y y)_{ij} \\
C_{ij}^{(k)} &= 1 + (\delta_x y)_{ij}^2 \\
D_{ij}^{(k)} &= -(\delta_y y)_{ij} \\
E_{ij}^{(k)} &= -(\delta_y y)_{ij}^2 \omega_{ij}
\end{align*} \tag{2.4.1.2}
\]
The $\delta_x y$ and $\delta_\psi y$ are approximated using 3-point central differences, namely,

\[
(\delta_x y)_{ij} = \frac{y_{i+1,j} - y_{i-1,j}}{2\Delta x}
\]

\[
(\delta_\psi y)_{ij} = \frac{y_{i,j+1} - y_{i,j-1}}{2\Delta \psi}
\]

Using 3-point central differences is acceptable for small $R_e$. However, as in the conventional stream function-vorticity formulation, the convective term in (2.4.1.1) may have to be upwind or backward differenced for larger $R_e$, i.e., upwind the vorticity term $(\delta_\psi \omega)_{ij}$ in the expression $R_e D_{ij}(\delta_x \omega)_{ij}$, namely

\[
(\delta_x \omega)_{ij} = \frac{\omega_{ij} - \omega_{i-1,j}}{\Delta x} \quad \text{(2-point)}
\]

\[
(\delta_\psi \omega)_{ij} = \frac{3 \omega_{ij} - 4 \omega_{i-1,j} + \omega_{i-2,j}}{2\Delta x} \quad \text{(3-point)}
\]

Numerical instabilities of explicit finite difference methods can be simply related to the familiar concepts of static and dynamic instabilities. Although not a consideration in this thesis, dynamic instabilities are caused by too large a time step. Static instabilities result from the form of the finite difference equation. For non-oscillatory solutions, limitations on the maximum Reynolds number, $R_e$, based on the finite difference cell size or the characteristic length $\Delta x$ or $\Delta \psi$, called the cell Reynolds number or Peclet number, are necessary. Oscillatory solutions may occur if the cell $R_e$ is too large. Whether they actually do occur depends on the flow geometry and imposition of boundary conditions.
However, a first-order accurate method, using upwind differences for the convective term in (2.4.1.1), although feasible in principle, is not recommended, insofar as the effective \( R_e \) of the numerical solution is lowered by the numerical viscosity introduced by the first-order accurate upwind differences. This point will be discussed more fully in Chapter III.

### 2.4.2 Solution Procedure

By using 3-point central difference approximations, equation (2.4.1.1) can be expressed as

\[
\frac{A_{ij}^{(0)}}{\Delta x^2} (\phi_{i-1,j} - 2\phi_{ij} + \phi_{i+1,j})^{(n+1)} + \frac{B_{ij}^{(0)}}{4\Delta x \Delta \psi} (\phi_{i+1,j+1} + \phi_{i-1,j-1})^{(n+1)}
\]

\[- \phi_{i+1,j-1} - \phi_{i-1,j+1})^{(n+1)} + \frac{C_{ij}^{(0)}}{\Delta \psi^2} (\phi_{ij-1} - 2\phi_{ij} + \phi_{ij+1})^{(n+1)}
\]

\[ + \frac{\alpha R_e D_{ij}^{(0)}}{2\Delta x} (\phi_{i+1,j} - \phi_{i-1,j})^{(n+1)} + \frac{E_{ij}^{(0)}}{2\Delta \psi} (\phi_{i,j+1} - \phi_{i,j-1})^{(n+1)} = 0
\]

By rearranging, we have

\[
B_{ij}^{(0)} \frac{\Delta x}{4 \Delta \psi} \phi_{i-1,j-1}^{(n+1)} + \left[ A_{ij}^{(0)} - \frac{\alpha R_e D_{ij}^{(0)} \Delta x}{2} \right] \phi_{i-1,j}^{(n+1)} - B_{ij}^{(0)} \frac{\Delta x}{4 \Delta \psi} \phi_{i-1,j+1}^{(n+1)}
\]

\[ + \left[ C_{ij}^{(0)} \frac{\Delta x^2}{\Delta \psi^2} - E_{ij}^{(0)} \frac{\Delta x^2}{2 \Delta \psi} \right] \phi_{ij-1}^{(n+1)} - 2 \left[ A_{ij}^{(0)} + C_{ij}^{(0)} \frac{\Delta x^2}{\Delta \psi^2} \right] \phi_{ij}^{(n+1)}
\]

\[ + \left[ C_{ij}^{(0)} \frac{\Delta x^2}{\Delta \psi^2} + E_{ij}^{(0)} \frac{\Delta x^2}{2 \Delta \psi} \right] \phi_{ij+1}^{(n+1)} - B_{ij}^{(0)} \frac{\Delta x}{4 \Delta \psi} \phi_{i+1,j-1}^{(n+1)} - 0
\]

(2.4.2.1a)
which can be abbreviated as follows, where \( \phi \) is at \((n+1)\)

\[
\begin{align*}
    b_{ij} \phi_{i-1,j-1} &+ (a_{ij} - d_{ij}) \phi_{i-1,j} - b_{ij} \phi_{i-1,j+1} \\
    + (c_{ij} - e_{ij}) \phi_{i,j-1} &- 2(a_{ij} + c_{ij}) \phi_{ij} + (c_{ij} + e_{ij}) \phi_{i,j+1} \\
    - b_{ij} \phi_{i+1,j-1} &+ (a_{ij} + d_{ij}) \phi_{i+1,j} + b_{ij} \phi_{i+1,j+1} = 0
\end{align*}
\]

(2.4.2.2a)

where

\[
\begin{align*}
    a_{ij} &= A_{ij}^{(1)} \\
    b_{ij} &= B_{ij}^{(1)} \frac{\Delta x}{\Delta \psi} \\
    c_{ij} &= C_{ij}^{(1)} \frac{\Delta x^2}{\Delta \psi^2} \\
    d_{ij} &= \alpha R_{ij} \frac{\Delta x}{2} \\
    e_{ij} &= E_{ij}^{(1)} \frac{\Delta x^2}{2 \Delta \psi}
\end{align*}
\]

(2.4.2.3a)

For a rectangular domain meshed with an \( IX \times JX \) grid, with known boundary values \( \phi_{ij} \) on the four boundaries where \( i = 1 \) or \( IX \), and \( j = 1 \) or \( JX \), the finite difference equations (2.4.2.2a) can be expressed in block tridiagonal matrix equation form as

\[
\begin{bmatrix}
    B_2 & C_2 \\
    A_3 & B_3 & C_3 \\
    & A_4 & B_4 & \cdots \\
    & & \cdots & \cdots \\
    & & & A_{12} & B_{12} & C_{12} \\
    & & & & A_{11} & B_{11}
\end{bmatrix} \begin{bmatrix}
    \phi_2 \\
    \phi_3 \\
    \phi_4 \\
    \vdots \\
    \phi_{12} \\
    \phi_{11}
\end{bmatrix}^{(n+1)} = \begin{bmatrix}
    R\text{H}S_2 \\
    R\text{H}S_3 \\
    R\text{H}S_4 \\
    \vdots \\
    R\text{H}S_{12} \\
    R\text{H}S_{11}
\end{bmatrix}
\]

(2.4.2.4)
where,

\[ \varphi_i^{(n+1)} = [\varphi_{i2}^{(n+1)}, \varphi_{i3}^{(n+1)}, \ldots, \varphi_{i1}^{(n+1)}]^T \]

is the solution vector along grid line i for \( i = 2, 3, \ldots, I_1 \). Here \( I_1 = I_X - 1, I_2 = I_X - 2 \), \( J_1 = J_X - 1 \) and \( J_2 = J_X - 2 \).

Matrices \( A_i, B_i, \) and \( C_i \) are \( J_2 \times J_2 \) tridiagonal matrices which can be expressed as

\[
A_i = \text{trid} \begin{bmatrix}
\text{for } 3 \leq i \leq I_1 & 3 \leq j \leq J_1 & 2 \leq j \leq J_1 & 2 \leq j \leq J_2 \\
b_{ij} & a_{ij} - d_{ij} & -b_{ij} & \end{bmatrix}
\]

\[
B_i = \text{trid} \begin{bmatrix}
\text{for } 2 \leq i \leq I_1 & 3 \leq j \leq J_1 & 2 \leq j \leq J_1 & 2 \leq j \leq J_2 \\
c_{ij} - e_{ij} & -2(a_{ij} + c_{ij}) & c_{ij} + e_{ij} & \end{bmatrix}
\]

\[
C_i = \text{trid} \begin{bmatrix}
\text{for } 2 \leq i \leq I_2 & 3 \leq j \leq J_1 & 2 \leq j \leq J_1 & 2 \leq j \leq J_2 \\
-b_{ij} & a_{ij} + d_{ij} & b_{ij} & \end{bmatrix}
\]

Abbreviating the above, we have

\[
A_i = \text{trid} (b_{ij}, a_{ij} - d_{ij}, -b_{ij})
\]

\[
B_i = \text{trid} (c_{ij} - e_{ij}, -2(a_{ij} + c_{ij}), c_{ij} + e_{ij})
\]

\[
C_i = \text{trid} (-b_{ij}, a_{ij} + d_{ij}, b_{ij})
\]

For example, if \( i = 3 \), we have
\[
A_3 = \begin{bmatrix}
  a_{32} - d_{32} & -b_{32} \\
  b_{33} & a_{33} - d_{33} & -b_{33} \\
  & & & \ddots & \ddots \\
  & & & & b_{3,J1} & a_{3,J1} - d_{3,J1}
\end{bmatrix}
\]

\[
B_3 = \begin{bmatrix}
  -2(a_{32} + c_{32}) & c_{32} + e_{32} \\
  c_{33} - e_{33} & -2(a_{33} + c_{33}) & c_{33} + e_{33} \\
  & & & \ddots & \ddots \\
  & & & & c_{3,J1} - e_{3,J1} & -2(a_{3,J1} + c_{3,J1})
\end{bmatrix}
\]

\[
C_3 = \begin{bmatrix}
  a_{32} + d_{32} & b_{32} \\
  -b_{33} & a_{33} + d_{33} & b_{33} \\
  & & & \ddots & \ddots \\
  & & & & -b_{3,J1} & a_{3,J1} + d_{3,J1}
\end{bmatrix}
\]

\[R\mathbf{H}_i\] is a \(I2 \times 1\) column vector which contains the boundary conditions and can be expressed as follows:

For \(i = 2\) or \(i = I1\), and \(j = 3, 4, \ldots, I2\)

\[
R\mathbf{H}_i = \begin{bmatrix}
  \gamma b_{i2} \phi_{i+1,1} - (a_{i2} + \gamma d_{i2}) \phi_{i+1,2} - \gamma b_{i2} \phi_{i+1,3} \\
  -\gamma b_{i2} \phi_{i+1,1} - (c_{i2} - e_{i2}) \phi_{i1} \\
  \vdots \\
  \gamma b_{i,j} \phi_{i+1,j-1} - (a_{ij} + \gamma d_{ij}) \phi_{i+1,j} - \gamma b_{i,j} \phi_{i+1,j+1} \\
  \vdots \\
  \gamma b_{i,I1} \phi_{i+1,I2} - (a_{i,I1} + \gamma d_{i,I1}) \phi_{i+1,I1} - \gamma b_{i,I1} \phi_{i+1,I1} \\
  + \gamma b_{i,I1} \phi_{i-1,I1} - (c_{i,I1} + e_{i,I1}) \phi_{i,JX}
\end{bmatrix}
\]
where

\[ \gamma = \begin{cases} 
-1 & \text{if } i=2 \\
1 & \text{if } i=11 
\end{cases} \]

Or more explicitly, if \( i = 2 \), then, for \( j = 3,4,...,J2 \)

\[
\text{RHS}_2 = \begin{bmatrix}
    b_{22}(\phi_{13}-\phi_{11})-(a_{22}-d_{22})\phi_{12} \\
    +b_{23}\phi_{31}-(c_{22}-e_{22})\phi_{21} \\
    \vdots \\
    b_{2J_2}\phi_{J1J2}-(a_{2J2}-d_{2J2})\phi_{J11} \\
    -b_{2J1}\phi_{3JX}-(c_{2J1}+e_{2J1})\phi_{2JX}
\end{bmatrix}
\]

If \( i = 11 \),

\[
\text{RHS}_{11} = \begin{bmatrix}
    b_{11,2}(\phi_{DX,1}-\phi_{DX,3})-(a_{11,2}+d_{11,2})\phi_{DX,2} \\
    -b_{11,2}\phi_{21,1}-(c_{11,2}-e_{11,2})\phi_{11,1} \\
    \vdots \\
    b_{11,J_2}(\phi_{DX,J1}-\phi_{DX,J2})-(a_{11,J2}+d_{11,J2})\phi_{DX,J1} \\
    \vdots \\
    b_{11,J1}(\phi_{DX,J2}-\phi_{DX,JX})-(a_{11,J2}+d_{11,J1})\phi_{DX,J1} \\
    +b_{11,J1}\phi_{3JX}-(c_{11,J1}+e_{11,J1})\phi_{2JX}
\end{bmatrix}
\]
For \(3 \leq i \leq I2\),

\[
RHS_i = \begin{bmatrix}
-b_2\phi_{i-1,i}-(c_2-e_2)\phi_{i1}+b_2\phi_{i+1,i} \\
0 \\
\cdot \\
\cdot \\
\cdot \\
0 \\
b_i\phi_{i-1,i+1}-(c_i+e_i)\phi_{i,i+1}-b_{i1}\phi_{i+1,i+1}
\end{bmatrix}
\]

The matrix equation is solved by an iterative method. The basic procedure is to obtain the solution using the method of successive line over-relaxation (SLOR), sweeping from left to right through the grid. The matrix equation (2.4.2.4) can be written vertically line by line so that on line \(i\), we have

\[
B_i\phi_i^{(n+1)} - RHS_i - A_i\phi_i^{(n+1)} - C_i\phi_i^{(n)}
\]

In pointwise form, this can be written as

\[
[B_i\phi_i^{(n+1)}]_j = RHS_y + b_y(\phi_{i-1,j+1}^{(n+1)} - \phi_{i-1,j-1}^{(n+1)}) - (a_y - d_y)\phi_{i-1,j}^{(n+1)} + b_y(\phi_{i-1,j-1}^{(n)} - \phi_{i-1,j+1}^{(n)}) - (a_y + d_y)\phi_{i-1,j}^{(n)}
\]

\[
- RHS_y + b_y(\phi_{i-1,j+1}^{(n+1)} - \phi_{i-1,j-1}^{(n+1)} + \phi_{i+1,j-1}^{(n)} - \phi_{i+1,j+1}^{(n)})
\]

\[
- (a_y - d_y)\phi_{i-1,j}^{(n+1)} - (a_y + d_y)\phi_{i+1,j}^{(n)}
\]

for all interior points, i.e., \(i = 3,4,\ldots,I2\) and \(j = 3,4,\ldots,I2\). Due to the boundary conditions along \(j = 1, j = JX, i = 1\) and \(i = IX\), we obtain the following.

If \(j = 2\),

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\[
[B_i \phi_i^{(n+1)}]_2 - RHS_{i2} = (a_{i2} - d_{i,2})\phi_{i-1,2}^{(n+1)} - (a_{i2} + d_{i,2})\phi_{i+1,2}^{(n)} + b_{i2}(\phi_{i-1,3}^{(n+1)} - \phi_{i+1,3}^{(n)})
\]

If \( j = J1 \),

\[
[B_i \phi_i^{(n+1)}]_{J1} - RHS_{i,J1} = (a_{i,J1} - d_{i,J1})\phi_{i-1,J1}^{(n+1)} - (a_{i,J1} + d_{i,J1})\phi_{i+1,J1}^{(n)} - b_{i,J1}(\phi_{i-1,J2}^{(n+1)} - \phi_{i+1,J2}^{(n)})
\]

For \( i = 2 \), for all \( j \) we have

\[
[B_2 \phi_2^{(n+1)}]_j - RHS_{2,j} = b_{2j}(\phi_{3,j-1}^{(n)} - \phi_{3,j+1}^{(n)}) - (a_{2j} + d_{2j})\phi_{3,j}^{(n)}
\]

If however \( j = 2 \), this further reduces to

\[
[B_2 \phi_2^{(n+1)}]_2 - RHS_{22} = b_{22}\phi_{33}^{(n)} - (a_{22} + d_{22})\phi_{32}^{(n)}
\]

or, if \( j = J1 \),

\[
[B_2 \phi_2^{(n+1)}]_{J1} - RHS_{2,J1} = b_{2,J1}\phi_{3,J2}^{(n)} - (a_{2,J1} + d_{2,J1})\phi_{3,J1}^{(n)}
\]

For \( i = I1 \), we have, for all \( j \)

\[
[B_i \phi_i^{(n+1)}]_j - RHS_{i,j} = b_{i,j}(\phi_{i+1,j-1}^{(n+1)} - \phi_{i+1,j-1}^{(n+1)}) - (a_{i,j} - d_{i,j})\phi_{i+1,j}^{(n+1)}
\]

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If \( j = 2 \), this reduces to
\[
[B_H \phi_i^{(n+1)}]_2 = RHS_{ij} \quad b_{H,2} \phi_i^{(n+1)} - (a_{H,3} - d_{H,3}) \phi_i^{(n+1)}
\]

and, if \( j = J_1 \),
\[
[B_H \phi_i^{(n+1)}]_j = RHS_{ij} - b_{H,j} \phi_i^{(n+1)} - (a_{H,j} - d_{H,j}) \phi_i^{(n+1)}
\]

Having obtained \( \overline{\phi}_i^{(n+1)} \), and calling these values \( \overline{\phi}_i^{(n+1)} \), a modified or relaxed value \( \overline{\phi}_i^{(n+1)} \) can be obtained by using a relaxation factor \( \beta_\alpha \) and the following expression:
\[
\overline{\phi}_i^{(n+1)} = (1 - \beta_\alpha) \overline{\phi}_i^{(n)} + \beta_\alpha \overline{\phi}_i^{(n+1)}
\]

where \( 1 \leq \beta_\alpha < 2 \) if \( \alpha = 0 \) (i.e., \( \phi = y \), \( \beta_o = \beta_y \)) and \( 0 < \beta_\alpha \leq 1 \) if \( \alpha = 1 \) (i.e., \( \phi = \omega \), \( \beta_1 = \beta_\omega \)). That is, \( y \) is over-relaxed and \( \omega \) is under-relaxed.

2.4.3 Iterative Procedure

1. Set \( \alpha = 0 \) (therefore \( \phi = y \), \( k = n \)). Linearize by evaluating coefficients at previous iteration level \( (n) \) as indicated in equation (2.4.2.1a). Solve for \( y_i^{(n+1)} \) using SLOR.

2. Evaluate near-boundary values of the speed from
\[
(q_i^{(n+1)})^2 = \left[ 1 + \frac{(\delta_y y)^2}{(\delta_y y)^2} \right]_{ij}^{(n+1)} - \frac{C_i^{(n+1)}}{A_i^{(n+1)}}
\] (2.4.3.1)
for \( j = 2, 3, 4, \ldots \), (near lower boundary) and for \( j = J1, J2, J3, \ldots \), (near upper boundary) as needed for evaluation of \( \omega \) at the boundaries (cf. Section 2.4.5).

3. Evaluate boundary values of the vorticities \( \omega_{i} \) and \( \omega_{i,JX} \) from appropriate formulas. Call these values \( \omega_{i}^{(n+1)} \) and \( \omega_{i,JX}^{(n+1)} \). Obtain under-relaxed boundary values \( \omega_{i}^{(n+1)} \) and \( \omega_{i,JX}^{(n+1)} \) by using a smoothing or damping parameter \( \delta \) (see computational preliminaries in next section for choice of \( \delta \)):

\[
\omega_{ij}^{(n+1)} = (1-\delta)\omega_{ij}^{(n)} + \delta \omega_{ij}^{(n+1)}, \quad 0 < \delta \leq 1, \quad j=1 \text{ and } JX.
\]

4. Set \( \alpha = 1 \) (therefore \( \phi = \omega, \ k = n+1 \)). Linearize by evaluating coefficients at iteration level \((n+1)\) for \( y_{ij} \)'s and previous iteration level \((n)\) for \( \omega_{ij} \)'s and approximate the second order derivatives at level \((n+1)\). Solve for \( \omega_{ij}^{(n+1)} \) using SLOR.

5. Set \( n = n+1 \) and repeat steps 1 to 4 until some specified convergence criteria is met, e.g., the iterations could be stopped when, for all \((i,j)\), the maximum norm

\[
\| \omega \| = \max_{ij} |\omega_{ij}^{(n+1)} - \omega_{ij}^{(n)}| < \varepsilon
\]

where \( \varepsilon \) is user-specified, say \( \varepsilon = 0.2 \times 10^{-4} \). Using the above convergence criteria guarantees [II.1] that for all \((i,j)\),

\[
\| y \| = \max_{ij} |y_{ij}^{(n+1)} - y_{ij}^{(n)}| < \varepsilon
\]

In some applications this convergence criteria may prove unsatisfactory. In that event, the iterations could be stopped when, for all \((i,j)\), the maximum norms
\[ \| y \| - \max \left| y_{ij}^{(n+1)} - y_{ij}^{(n)} \right| < \varepsilon_y \]
\[ \| \omega \| - \max \left| \omega_{ij}^{(n+1)} - \omega_{ij}^{(n)} \right| < \varepsilon_{\omega} \]

where \( \varepsilon_y \) and \( \varepsilon_{\omega} \) are user-specified, say \( \varepsilon_y = 0.2 \times 10^{-3} \) and \( \varepsilon_{\omega} = 0.2 \times 10^{-2} \). The iterations start with some initial approximation \( \omega_{ij}^{(0)} \) of the vorticity with \( n = 0 \). If no such approximation exists, set \( \omega_{ij}^{(0)} = 0 \).

Equations for \( y \) and \( \omega \) are solved at interior points only, i.e., \( 2 \leq i \leq I1 = IX - 1 \), \( 2 \leq j \leq J1 = JX - 1 \). Hence, the fact that the Jacobian becomes infinite along the solid boundaries does not create a problem because we do not apply the partial differential equations on these boundaries, i.e., \( u - \frac{1}{y_{\psi}} = 0 \), \( v - \frac{y_x}{y_{\psi}} = 0 \), implying \( y_{\psi} \to \infty \) on the boundaries, would make the Jacobian infinite. Also, it is well established by computational experience that the iterative convergence can be judged by examining only the vorticity on the wall.

2.4.4 Computational Preliminaries

A non-zero value of the damping parameter \( \delta \) is essential for the convergence of the numerical procedure [II.1]. For an estimate of \( \delta \), one must determine the growth factor \( \rho \) for the \( y_{ij} \) iterations. The value of \( \rho \) is estimated by using \( \delta = 0 \) in step 3 of the iterative procedure for a small number of iterations. Compute, for large \( n \)

\[ \rho \approx \frac{\| y_{ij}^{(n+1)} - y_{ij}^{(n)} \|}{\| y_{ij}^{(n)} - y_{ij}^{(n-1)} \|} , \]
The norm used is the maximum norm \( \| y \| = \max_{ij} |y_{ij}| \). Let

\[
\mu = \frac{\rho - 1}{\rho + 1}
\]

For convergence of the \( y_{ij} \) iterations, \( \delta \) should be chosen such that \( \mu \leq \delta \leq 1 \) and a near optimal value of \( \delta \) is given by

\[
\delta_{opt} = \frac{\rho}{\rho + 2}
\]

2.4.5 Boundary Conditions for Vorticity

The values of vorticity on the solid boundaries are generally unknown and must be obtained as part of the overall solution. As in the conventional \((\psi, \omega)\) formulation, these values of vorticity must be approximated along any solid boundaries.

By definition,

\[
\omega(x,y) = v_x - u_y \quad (2.4.5.1)
\]

Also, by definition of the stream function

\[
\begin{align*}
\psi &= \psi_y \\
v &= -\psi_x
\end{align*} \quad (2.4.5.2)
\]

Since \( u = u(x,y) \), \( v = v(x,y) \), and under the von Mises transformation, \( y = y(x,\psi) \), we get \( u = u(x,\psi) \), \( v = v(x,\psi) \). Hence,

\[
\begin{align*}
v_x &= v_x x_x + v_x \psi_x = v_x - vv_y \\
u_y &= u_x x_y + u_x \psi_y = uu_y
\end{align*} \quad (2.4.5.3)
\]
using (2.4.5.2). Substituting (2.4.5.3) into (2.4.5.1), we get

$$\omega = \omega(x, \psi) = v_x - vv_\psi - uu_\psi$$  

(2.4.5.4)

In [II.1] several solid boundary approximations are suggested for vorticity. However, expanding any of these suggested schemes in a Taylor series about the point \((i,1)\) \(i = 1, 2, 3, \ldots\) and using the no-slip condition on this boundary, namely \(u_{\text{in}} = v_{\text{in}} = 0\), equation (2.4.5.4) yields \(\omega_{\text{in}} = 0\) for all \(i\). Obviously, this is no good for our purposes. The problem seems to be that even though \(u = v = 0\) on the boundary, \(uu_\psi\) and \(vv_\psi\) do not tend to zero. In fact, from (2.4.5.3), \(uu_\psi \to 0\) would imply \(u_\psi \to 0\), which is obviously incorrect for physical reasons. Hence, an alternate approach is required.

Now, \(q^2 = u^2 + v^2\), and therefore \((q^2)_\psi = 2uu_\psi + 2vv_\psi\) or

$$\frac{1}{2}(q^2)_\psi = uu_\psi + vv_\psi$$  

(2.4.5.5)

Substituting (2.4.5.5) into (2.4.5.4) we get

$$\omega = \omega(x, \psi) = v_x - \frac{1}{2}(q^2)_\psi$$  

(2.4.5.6a)

Along a solid boundary the stream function \(\psi = \psi_b = \text{constant}\). In most of our applications, solid boundaries occur when \(i = 1\) or IX for all \(j\) and/or \(j = 1\) or JX for all \(i\). These usually correspond to curves where, for convenience, we take \(\psi = 0\) or \(\psi = \psi_{\text{MAX}}\).

The no-slip viscous boundary condition on solid boundaries, namely \(u = v = 0\) must be imposed. In particular, this implies that \(v_x = 0\) on \(\psi = \psi_b = \text{constant}\). Thus, (2.4.5.6a) reduces to

$$\omega(x, \psi) = -\frac{1}{2} \frac{\partial q^2}{\partial \psi} \bigg|_{\psi=\psi_b}$$  

(2.4.5.6b)
and clearly \( \frac{\partial q^2}{\partial \psi} \neq 0 \) at the boundary \( \psi = \psi_b \).

Hence, the following formulas for vorticity can be obtained by Taylor series expansions using the no-slip boundary condition on solid boundaries, namely \( q_{i_n}^2 = q_{i_nX}^2 = 0 \) for a bottom or top boundary, to obtain \( \omega_{i_n} \) and \( \omega_{i_nX} \), respectively (\( i = 2,3,4,\ldots \)).

For the bottom boundary \( \psi = 0 \) (\( j = 1 \)), one of the following approximate relations can be used (cf. Appendix F):

\[
0(\Delta \psi): \quad \omega_{i_n} \approx -\frac{1}{2\Delta \psi} q_{i_n}^2 \tag{2.4.5.7a}
\]

\[
\omega_{i_n} \approx -\frac{1}{2\Delta \psi} (q_{i_5}^2 - q_{i_3}^2) \tag{2.4.5.7b}
\]

\[
\omega_{i_n} \approx -\frac{1}{2\Delta \psi} (q_{i_5}^2 - q_{i_3}^2) \tag{2.4.5.7c}
\]

\[
0(\Delta \psi^2): \quad \omega_{i_n} \approx \frac{1}{4\Delta \psi} (q_{i_4}^2 - 4q_{i_2}^2) \tag{2.4.5.7d}
\]

\[
\omega_{i_n} \approx \frac{1}{12\Delta \psi} (4q_{i_4}^2 - 9q_{i_3}^2) \tag{2.4.5.7e}
\]

\[
\omega_{i_n} \approx \frac{1}{4\Delta \psi} (3q_{i_4}^2 - 8q_{i_2}^2 + 5q_{i_1}^2) \tag{2.4.5.7f}
\]

\[
0(\Delta \psi^3): \quad \omega_{i_n} \approx -\frac{1}{12\Delta \psi} (2q_{i_4}^2 - 9q_{i_2}^2 + 18q_{i_1}^2) \tag{2.4.5.7g}
\]

\[
\omega_{i_n} \approx -\frac{1}{12\Delta \psi} (11q_{i_4}^2 - 42q_{i_2}^2 + 57q_{i_3}^2 - 26q_{i_5}^2) \tag{2.4.5.7h}
\]

\[
0(\Delta \psi^4): \quad \omega_{i_n} \approx -\frac{1}{72\Delta \psi} (q_{i_2}^2 + 8q_{i_4}^2 - 48q_{i_3}^2 + 104q_{i_5}^2) \tag{2.4.5.7i}
\]

Similarly, for a top boundary \( \psi = \psi_{MAX} \) (\( j = JX \)), one of the following approximate relations can be applied:
\[ 0(\Delta \psi): \quad \omega_{i,jx} = \frac{1}{2\Delta \psi} q_{i,j1}^2 \]  \hspace{1cm} (2.4.5.8a)
\[ \omega_{i,jx} = \frac{1}{2\Delta \psi} (q_{i,j2}^2 - q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8b)
\[ \omega_{i,jx} = \frac{1}{2\Delta \psi} (q_{i,j3}^2 - q_{i,j2}^2) \]  \hspace{1cm} (2.4.5.8c)

\[ 0(\Delta \psi^2): \quad \omega_{i,jx} = -\frac{1}{4\Delta \psi} (q_{i,j2}^2 - 4q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8d)
\[ \omega_{i,jx} = -\frac{1}{12\Delta \psi} (4q_{i,j3}^2 - 9q_{i,j2}^2) \]  \hspace{1cm} (2.4.5.8e)
\[ \omega_{i,jx} = -\frac{1}{4\Delta \psi} (3q_{i,j3}^2 - 8q_{i,j2}^2 + 5q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8f)

\[ 0(\Delta \psi^3): \quad \omega_{i,jx} = \frac{1}{12\Delta \psi} (2q_{i,j3}^2 - 9q_{i,j2}^2 + 18q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8g)
\[ \omega_{i,jx} = \frac{1}{12\Delta \psi} (11q_{i,j4}^2 - 42q_{i,j3}^2 + 57q_{i,j2}^2 - 26q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8h)

\[ 0(\Delta \psi^4): \quad \omega_{i,jx} = \frac{1}{72\Delta \psi} (q_{i,j4}^2 + 8q_{i,j3}^2 - 48q_{i,j2}^2 + 104q_{i,j1}^2) \]  \hspace{1cm} (2.4.5.8i)

Note that the positive and negative signs in equations (2.4.5.7) and (2.4.5.8) are correct as indicated.

At first glance, some of these approximations may appear wrong or inconsistent. For example, equations (2.4.5.7b) and (2.4.5.7d) involve the same points (i,2) and (i,3), but (2.4.5.7b) was obtained by writing
\[ \omega_{ii} = a q_{ij}^2 + b q_{kl}^2 \]

while (2.4.5.7d) was obtained by writing

\[ \omega_{ii} = a q_{ij}^2 + b q_{kl}^2 + c q_{il}^2 \]

and, after finding the constants \(a, b, c\), using the no-slip boundary condition to set \(q_{ii}^2 = 0\). Derivation of equations (2.4.5.7b) and (2.4.5.7d) is contained in Appendix F to demonstrate more clearly the significance of this point. Obviously, (2.4.5.7d) is the correct result to use since it employs the non-slip condition while (2.4.5.7b) does not.

The formulas (2.4.5.7b, c, f, h) are symmetric in the sense that we do not use the term \(q_{ii}^2\) in the derivation, whereas the formulas (2.4.5.7a, d, e, g, i) are non-symmetric in the sense that we take \(q_{ii}^2 = 0\) after calculating the constant coefficients. Further, note that the value of the constants in the symmetric formulas sums to zero as they should, but in the non-symmetric formulas they do not; however, the coefficient of \(q_{ii}^2\) can be easily calculated. For example, in formula (2.4.5.7d), summing the constants to zero implies that the constant coefficient for \(q_{ii}^2\) must be +3, as noted in Appendix F. Although not a consideration for the problems in this thesis, if the values of the \(q_j^2\) \((j = 1, 2, 3, \ldots)\) happen to be close to 1, the resulting value for \(\omega_{ii}\) could be in error when using the symmetric formulas. More investigation of this situation is required in future research.

Some preliminary testing of the boundary conditions for vorticity \(\omega_{ii}\) was undertaken for a simple test problem to try to determine the most appropriate formula to use. Knowing exact results for \(\omega_{ii}\) at \(x = 0\) (i.e., \(i = 1\)) the various formulas were tested one-by-one to determine the most accurate. Of all the formulas listed, equation
(2.4.5.7g) (and its corresponding formula (2.4.5.8g) for \( \omega_{1,xx} \)) came closest to the exact value. These equations will be used as applicable, for the lower and upper boundary values of vorticity, respectively.

However, the choice of boundary condition also affects numerical stability and the rate of convergence of the iterative process. Another possible subject for future research will be the systematic comparison of the solutions obtained using the different expressions for vorticity.

2.4.6 Clustered Grids

To achieve clustering of the grids in regions of high gradients, one-dimensional stretching functions are employed.

The flow equations with appropriate boundary conditions have been presented in unstretched coordinates. In order to formulate the equations in the stretched coordinates, general transformations can be introduced, defined by

\[
x = x(\xi) \\
\psi = \psi(\eta)
\] (2.4.6.1)

The benefit of these transformations is that they can provide us with a dense mesh in the vicinity of any singularity and allow us to pack more points near an axis or any solid boundary.

Equation (2.4.1.1) and coefficients (2.4.1.2) are transformed from unstretched coordinates \((x, \psi)\) to stretched coordinates \((\xi, \eta)\) for later use. Initially, only the second transformation of equation (2.4.6.1) will be used and the required equations derived, then both transformations will be applied. Our reason for doing this is dictated by the nature of the particular problems to be solved, for example, some may require only the \(\psi\)
coordinate to be stretched, while some will require stretching both the x and \( \psi \) coordinate. Details and the resulting equations can be found in Appendix G.

2.4.7 Large Reynolds Number

The central differencing scheme employed to obtain (2.4.2.1a) works well for low \( R_e \). For larger \( R_e \), it becomes necessary to upwind or backward difference the convective term \((\delta_x\omega)_j\) in (2.4.1.1). In this case \( \alpha = 1, \phi = \omega(x, \psi) \) and equation (2.4.1.1) can be expressed as

\[
\frac{A_{ij}^{(a)}}{\Delta x^2} (\omega_{i-1,j} - 2\omega_{ij} + \omega_{i+1,j})^{(a+1)} + \frac{B_{ij}^{(a)}}{4\Delta x \Delta \psi} (\omega_{i+1,j+1} - \omega_{i-1,j-1})^{(a+1)}
\]

\[
-\omega_{i+1,j-1} - (\omega_{i,j+1} - \omega_{i+1,j})^{(a+1)} + \frac{C_{ij}^{(a)}}{\Delta \psi^2} (\omega_{i,j-1} - 2\omega_{ij} + \omega_{i,j+1})^{(a+1)}
\]

\[
+ \frac{R_e D_{ij}^{(a)}}{\Delta x} (\omega_{ij} - \omega_{i-1,j})^{(a+1)} + \frac{E_{ij}^{(a)}}{2\Delta \psi} (\omega_{i-1,j} - \omega_{i,j-1})^{(a+1)} = 0
\]

By rearranging, we have

\[
B_{ij}^{(a)} \frac{\Delta x}{4\Delta \psi} \omega_{i-1,j-1}^{(a+1)} + \left( A_{ij}^{(a)} - R_e D_{ij}^{(a)} \Delta x \right) \omega_{i-1,j}^{(a+1)} - B_{ij}^{(a)} \frac{\Delta x}{4\Delta \psi} \omega_{i-1,j+1}^{(a+1)}
\]

\[
+ \left[ C_{ij}^{(a)} \frac{\Delta x^2}{\Delta \psi^2} - E_{ij}^{(a)} \frac{\Delta x^2}{2\Delta \psi} \right] \omega_{i,j-1}^{(a+1)} - 2 \left[ A_{ij}^{(a)} + C_{ij}^{(a)} \frac{\Delta x^2}{\Delta \psi^2} + R_e D_{ij}^{(a)} \Delta x \right] \omega_{ij}^{(a+1)}
\]

\[
+ \left[ C_{ij}^{(a)} \frac{\Delta x^2}{\Delta \psi^2} + E_{ij}^{(a)} \frac{\Delta x^2}{2\Delta \psi} \right] \omega_{i+1,j-1}^{(a+1)} - B_{ij}^{(a)} \frac{\Delta x}{4\Delta \psi} \omega_{i+1,j+1}^{(a+1)}
\]

\[
+ A_{ij}^{(a)} \omega_{i+1,j}^{(a+1)} + B_{ij}^{(a)} \frac{\Delta x}{4\Delta \psi} \omega_{i+1,j+1}^{(a+1)} = 0
\]

which can be abbreviated as
\[ b_{ij} \omega_{i-1,j+1} + (a_{ij} - d_{ij}^*) \omega_{i,j} - b_{ij} \omega_{i+1,j+1} + (c_{ij} - e_{ij}) \omega_{i+1,j} - 2(a_{ij} + c_{ij} + d_{ij}^*) \omega_{i,j} + (c_{ij} + e_{ij}) \omega_{i+1,j+1} = 0 \tag{2.4.2.2b} \]

where,

\[ a_{ij} = A_{ij}^{(o)} \]
\[ b_{ij} = B_{ij}^{(o)} \frac{\Delta x}{4\Delta \psi} \]
\[ c_{ij} = C_{ij}^{(o)} \frac{\Delta x^2}{\Delta \psi^2} \tag{2.4.2.3b} \]
\[ d_{ij}^* = R_{ij} D_{ij}^{(o)} \Delta x \]
\[ e_{ij} = E_{ij}^{(o)} \frac{\Delta x^2}{2\Delta \psi} \]

Corresponding changes have to be made to the various tridiagonal matrices. Comparing equations (2.4.2.2a) and (2.4.2.2b), it is seen that the upwind differencing of the convective term only changes the coefficients of \( \omega_{i-1,j} \), \( \omega_{ij} \) and \( \omega_{i+1,j} \). The * in \( d_{ij}^* \) is used so as not to cause confusion with the expression for \( d_{ij} \) stated previously in (2.4.2.3a).
CHAPTER III

PLANE DIVERGING VISCOUS LAMINAR CHANNEL FLOW OF COMPLEX GEOMETRY

3.1 INTRODUCTION

In 1982, the International Association for Hydraulic Research (IAHR) Working Group on Refined Modelling of Flows devoted the Fifth IAHR Meeting to a specific subject - to assess the capabilities of various numerical simulation methods to deal with laminar flows in "complex geometries". Here "complex geometries" means flow domains that do not coincide with coordinate axes in some simple coordinate system such as Cartesian or polar.

Why study such flows? Steady-state viscous flow in two-dimensional channels with arbitrary wall contours are representative of flush inlet geometries [III.1]. The water inlets for ships powered by water pumps are often mounted flush to the hull below the water line. The calculation of the flow in flush inlets is of interest for the prediction of the total drag of the ship and the performance of the pump. Flow separation can occur in the inlet, so that viscous effects cannot be ignored.

A single, well defined comparison test problem, namely the laminar flow in a channel with a smooth expansion, suggested by the work of Roache [III.2] on the scaling of Reynolds number in weakly separated (i.e., small recirculating region) channel flows, was chosen for testing various numerical methods. The purpose of the test problem was to evaluate the capabilities of various Navier-Stokes solvers and to highlight difficulties in the modelling of complex geometries. This problem has been used to test the present formulation. A comparison and discussion of the solutions obtained by the participants was reported by Napolitano and Orlandi [III.3]. As reported in [III.3], some of the
participants felt that the problem under investigation was 'too easy' and therefore not suitable to assess the capability of each code to compute flows in complex geometries. As the author of this thesis and others discovered, and as the reader will verify after analysis of the results obtained, such an opinion was premature.

3.2 TEST PROBLEM

3.2.1 Problem Specification

The geometry is a diverging channel with length depending on the Reynolds number $R_e$, i.e., the length of the channel is scaled proportionally to $R_e$ so that the channel becomes longer and straighter as $R_e$ increases, i.e., $R_e = R_{ce}$ is a geometrical constant which determines the steepness of the curved wall of the expansion (see Diagram 3.1.1). For $R_e > 1$, quasi-self-similar flow conditions and solutions can be obtained [III.2] by having the channel length $x$ increase proportional to $R_e$ so solutions become self-similar in the scaled longitudinal variable $x_{out} = R_{ce}/3$ ($x_{out} =$ outlet of channel). This scaling is necessary to keep the separated flow region within the computational mesh. Weakly separated laminar two-dimensional incompressible channel flows display a self-similar solution.

Two flows were computed corresponding to relatively small values of $R_e$, i.e., $R_e = R_{ce} = 10$ and $R_e = R_{ce} = 100$. $R_e = 10$ was chosen because of its highly distorted geometry. $R_e = 100$ was chosen to assess the dependence of the convergence rate on $R_e$. A 21 x 21 finite difference mesh was prescribed. Computed results for the wall were obtained for equally spaced $x/x_{out}$ locations.

Because the numerical results from these two problems should not depend too much on the treatment of advection terms, an optional third case was suggested, namely $R_e = 100$ inside the $R_{ce} = 10$ channel. Even though this case would be characterized
by more significant advection, a case more strongly influenced by the modelling of non-linear terms, it was not considered simply because this third problem was not physically representative. As noted in [III.4], flow in a symmetric channel at high Reynolds number is not unique. In the case of laminar flow, the existence of asymmetric solutions implies non-uniqueness of solutions of the Navier-Stokes equations. Also, the symmetric solution modelled by the test problem is, in fact, unstable at high Reynolds number. Thus, the third test problem, although it can be studied computationally, is not physically realistic.

3.2.2 Boundary Conditions in the Physical Domain

The lower boundary (solid wall) coordinates of the channel are given analytically as

\[ y - y_s(x) = \frac{1}{2} \left[ \tanh \left( 2 - \frac{30x}{R_E} \right) - \tanh 2 \right], \quad 0 \leq x \leq x_{\text{out}} - \frac{R_{\text{ce}}}{3} \quad (3.2.2.1) \]

and no-slip conditions \( u = v = 0 \) are applied along this boundary. (See Appendix H for a discussion of this function).

The upper boundary (centreline or symmetry plane) is located at

\[ y - y_u(x) = 1.0, \quad 0 \leq x \leq x_{\text{out}} \]

Inlet boundary conditions are given in terms of the Cartesian velocity components \((u,v)\) as

\[ u = 3 \left( y - \frac{y^2}{2} \right) \quad \text{for } x = 0, \quad 0 \leq y \leq 1 \]

\[ v = 0 \]
i.e., equilibrium flow of the inlet with an imposed fully developed parabolic Poiseuille flow velocity distribution. The origin of the physical coordinates \((x, y)\) is on the lower wall at the inflow boundary, where the channel half height has been normalized to \(y = 1\). The maximum inflow velocity is \(u(0,1) = 3/2\). The length of the channel is \(x = x_{\text{out}} = R_{\text{sc}}/3\).

Several investigators used a simple linear transformation or stretching to map \((x, y)\) into a rectangular computational domain. We also obtain a rectangular region for the computational domain (Diagram 3.1.2), but the transformation to achieve this has been done on the original partial differential equations. Finally, it should be noted that Poiseuille flow implies constant area flow, which in the case of a diverging channel is a contradiction. This results in a particularly annoying difficulty, to be discussed later.

In terms of the stream function, the inlet conditions are

\[
\begin{align*}
\psi_y &= 3 \left( y - \frac{y^2}{2} \right) \\
\psi_x &= 0
\end{align*}
\]

Therefore,

\[
\psi = \psi(y) = \frac{1}{2} (3y^2 - y^3) \quad \text{for } x = 0, \ 0 \leq y \leq 1
\]

where we have taken \(\psi = 0\) at \(y = 0\).

The standard no-slip condition on solid walls (i.e., \(u = v = 0\)) is imposed at the wall \(0 \leq x \leq x_{\text{out}}\) and \(y = y_t(x)\). Symmetry is enforced at \(0 \leq x \leq x_{\text{out}}\) and \(y = y_t(x) = 1\).

The outlet section of the channel is located at
\[ x - x_{\text{out}} = \frac{R_{cc}}{3} \]

so that, at the lower boundary

\[ y - y_t(x_{\text{out}}) = \frac{1}{2} [ \tanh(-8) - \tanh2 ] = -0.98202 \]

The outlet boundary conditions are left somewhat arbitrary [III.2]. For parallel flow

\[ u_x = v_x = 0, \text{ for } x = x_{\text{out}} \text{ and } y_t \left( \frac{R_{cc}}{3} \right) \leq y \leq 1 . \]

This implies \( y_x = \omega_x = 0 \). See Appendix I for a further discussion of this point.

3.2.3 Boundary Conditions in the Computational Domain

Boundary conditions for \( y \):

- At \( x = 0, 0 \leq \psi \leq 1 \)

\[ \psi = \frac{1}{2}(3\psi^2 - \psi^3) \]

which can be solved explicitly for \( y \) as a function of \( \psi \) to give (see Appendix J)

\[ y = 2\cos \left( \theta + \frac{4\pi}{3} \right) + 1 \quad (3.2.3.1) \]

where

\[ \theta = \frac{1}{3} \arccos(1 - \psi) \]
\begin{itemize}
  \item At \( x = x_{\text{out}} \), \( 0 \leq \psi \leq 1 \)
    \[ y_x = 0 \]
  \item At \( \psi = 0 \), \( 0 \leq x \leq x_{\text{out}} \)
    \[ y = y_t(x) = \frac{1}{2} \left[ \tanh \left( 2 - \frac{30x}{R_c} \right) - \tanh 2 \right] \]
  \item At \( \psi = 1 \), \( 0 \leq x \leq x_{\text{out}} \)
    \[ y = y_u(x) = 1 \]
\end{itemize}

Boundary conditions for \( \omega \):

\begin{itemize}
  \item At \( x = 0 \), \( 0 \leq \psi \leq 1 \), by definition \( \omega = -\nabla^2 \psi \) and \( \psi = \psi(y) \) at the inlet, so that
    \[ \omega = -\psi_{yy} = -u_y = -3(1-y) \]  
    \hspace{1cm} (3.2.3.2)
  \item At \( x = x_{\text{out}} \), \( 0 \leq \psi \leq 1 \)
    \[ \omega_x = 0 \]
  \item At \( \psi = 0 \), \( 0 \leq x \leq x_{\text{out}} \)
    \[ \omega_3 = -\frac{1}{12\Delta \psi} \left( 2q_{11}^2 - 9q_{13}^2 + 18q_{12}^2 \right) \]
  \item At \( \psi = 1 \), \( 0 \leq x \leq x_{\text{out}} \)
    \[ \omega = 0 \] (symmetry condition)
\end{itemize}

Note: At \( x = x_{\text{out}} \), \( 0 \leq \psi \leq 1 \), 3-point backward or upward differencing has been used for \( y_x = \omega_x = 0 \), i.e., \( \frac{\phi_{1z,j} - 4\phi_{1z,j+1} + 3\phi_{1z,j-1}}{2\Delta x} = 0 \), \( \phi = \{ y \omega \} \).

\section{Vorticity Discontinuity at the Inlet}

This section could have been placed under the heading of discussion. However,
because of the significance of this issue, it was felt that a discussion of this topic should come earlier.

Unfortunately, the test problem had a defect [III.4]. As noted earlier, the validity of the inlet boundary conditions was questionable. Fully developed Poiseuille flow conditions have been prescribed for the inlet velocity profile in spite of the non-zero slope of the wall of the channel at $x = 0$. As a consequence of this inappropriate choice of inlet profile, there is a singularity at this point, which shows up either as a disturbed wall pressure distribution or as a discontinuity in vorticity on the channel wall at the inlet or origin. In fact, the choice of the outlet as the reference pressure point in a primitive-variable approach would also have caused difficulties because of the arbitrariness of the boundary conditions at the outlet that the stream function-vorticity approach could alleviate.

As noted in [III.4], the magnitude of the discontinuity in vorticity at the origin because of the nature of the boundary conditions is a function of the angle the wall makes with the horizontal at the inlet. The discontinuity only affects the local flow, which can be made plausible by considering the velocity close to the wall. The effect, however, is rather pronounced. At the inlet, the local flow is parallel to the x axis, but just inside the inlet it must be parallel to the channel wall, which is not parallel to the x axis. By constructing the local flow solution, Cliffe et al [III.4] proved that there is in fact a discontinuity in vorticity and obtained a value for the jump. For the case $Re = Re_{cc} = 10$, the value obtained was 0.9125 and for the case $Re = Re_{cc} = 100$, the value obtained was 0.0989. Both of these values are used in this thesis.

In addition to the results obtained for the jump in vorticity at the inlet, the inlet boundary condition for vorticity must be modified. At $x = 0$, $0 \leq \psi \leq 1$, we obtained
equation (3.2.3.2), namely, $\omega = -3(1-y)$. This was obtained by noting that $\psi = \psi(y)$ only at the inlet and therefore that $\omega = -\psi_{yy} = -u_y$. Equally, this could have been obtained by using:

$$\omega = v_x - u_y$$

and the fact that,

$$u = 3 \left( y - \frac{y^2}{2} \right) = u(y)$$
$$v = 0$$

at the inlet. Taking $v_x = 0$ gives

$$\omega = -u_y$$

Actually, this is incorrect since $v_x \neq 0$ unless the channel is perfectly straight for a couple of grid points to the right of the inlet (a method also tried without boundary condition modification). Hence, the boundary condition for $\omega$ at the inlet (equation (3.2.3.2)) is modified as follows:

At $x = 0$, $0 \leq \psi \leq 1$,

$$\omega = -3(1-y) + v_x$$

(3.2.4.1)

where $v_x$ is approximated by a one-sided or forward (downwind) difference formula, either 2-point or 3-point. (See Appendix K for a summary of the difference formulas used for $v_x$). Note that at $y = 0$ ($i=1$), from the no-slip boundary condition along the wall, $v = 0$ still. This implies $v_x = 0$. Hence at $x = 0$, $y = 0$ ($i=1$, $j=1$) we still have $\omega_{11} = -3$. However, for $i = 1$, $j = 2, 3, ..., J1$, $v_x \neq 0$.

### 3.2.5 Clustered Grid Functions

The mesh distribution employed in the calculations is very important to modelling
the separation region correctly. In particular, using a proper stretching in the direction normal to the wall could be essential. In this way a finer resolution is obtained in the region where a separation bubble is likely to develop.

Hence, it is desirable to concentrate grid lines close to the channel wall first, and second, to redistribute lines away from the inlet to the body of the domain [III.5]. The first objective can be achieved by choosing \( \psi \) in equation (2.4.6.1) to be

\[
\psi = \eta_{\text{MAX}} \left[ 1 - e^{-a(\eta_{\text{MAX}} - \eta)} \right]
\]

(3.2.5.1)

where \( \eta_{\text{MAX}} \) is the largest value of \( \eta \) in the new computational domain. In our present case \( \eta_{\text{MAX}} = 1.0 \). The larger the value chosen for the constant \( \alpha \), the more grid lines will be concentrated towards the channel wall. The second objective can be achieved by setting \( x \) in equation (2.4.6.1) to be

\[
x = \sinh [\beta(\xi - \xi_0)]
\]

(3.2.5.2)

where \( \xi_0 \) is the grid line around which concentration of the grid lines is desired and \( \beta \) is again a constant which determines the degree of concentration.

3.3 RESULTS AND DISCUSSIONS

3.3.1 Preliminaries

In the absence of an exact reference solution, the grid-independent results obtained by Cliffe et al in [III.4] have been taken as a benchmark, as recommended by Napolitano and Orlandi [III.3]. Cliffe et al used a finite element method in primitive variables, a Newton-Raphson linearization scheme and the frontal solution method for the resulting linear system. Such a solution has been used to compute the average percentage error \( \varepsilon_{\omega} \) defined according to the relationship
\[
\varepsilon_{u_w} = \frac{100}{19} \sum_{i=2}^{20} \left| \frac{\omega_{i} - \omega_{i, CJG}}{\omega_{i, CJG}} \right|
\]  
(3.3.1.1)

where \(\omega_i\) is the computed wall vorticity at 21 equally spaced \(\frac{x}{x_{out}}\) locations along the wall and the subscript CJG refers to the benchmark solution of Cliffe et al noted above. The values of the gridpoints \(x = 0\) and \(x = x_{out}\) have not been included in the definition of \(\varepsilon_{u_w}\) to reduce the influences of the singularity at the inlet and of the arbitrary outlet boundary conditions. \(\varepsilon_{u_w}\) has been defined so as to account mostly for the region around and inside the separation bubble as the results of this thesis will indicate. Since this is a rational and appropriate choice, \(\varepsilon_{u_w}\) is a good quantity to judge the accuracy of the solutions for the present flow case (more so than, say, pressure). If the method under consideration ignores the separation phenomenon completely or if the length and the position of the separation region are not computed very accurately, then \(\varepsilon_{u_w}\) will be very large because very large relative errors for \(\omega\) are probable near the separation and the reattachment points. One must appreciate the goal of trying to obtain accuracy away from thin boundary layers without actually resolving these boundary layers in detail.

### 3.3.2 Results

Results were obtained for the following problems:

1. Reynolds number: (i) \(R_c = R_{\infty} = 10\) and (ii) \(R_c = R_{\infty} = 100\) where \(\Delta x = \frac{R_{\infty}}{3}/20\). For example, for \(R_{\infty} = 10\), \(\Delta x = 0.167\). For \(R_{\infty} = 100\), \(\Delta x = 1.667\). \(\Delta \psi = 0.05\) in both cases.

2. Inlet corrections: (i) No inlet correction, (ii) the inlet correction used by Cliffe et al [III.4] and (iii) the modified boundary condition given by equation (3.2.4.1).
For Reynolds number \( R_e = R_{ce} = 10 \), Figures (3.3.2.1), (3.3.2.2) and (3.3.2.3) contain plots of the streamlines \( y_j, j=1 \) (the wall), \( j=2 \) (first streamline) to \( j = 10 \) for \( i = 1,2,\ldots,20 \).

Figure (3.3.2.4) is a plot of the wall vorticity \( \omega_i, i = 1,2,\ldots,20 \).

Figure (3.3.2.5) is a plot of the worst case wall vorticity values \( \omega_i, i = 1,2,\ldots,20 \) obtained by all the participants that reported in [III.3] against the results obtained in Figure (3.3.2.4) \( (R_e = R_{ce} = 10) \).

Table (3.3.2.1) contains a listing of the various parameters used to obtain the results.

Tables (3.3.2.2), (3.3.2.3) and (3.3.2.4) contain the error for each point of the wall vorticity and the total error. The number of iterations referred to in the tables is for one equation. Hence, the number of iterations for the system of equations is half the quoted value.

Table (3.3.2.5) contains a comparison of the separation and reattachment points and the relative errors where the vorticity \( \omega = 0 \) for various values of \( x \).

All the results were obtained using an IBM PC-compatible/Intel 80286 (AT) computer.

3.3.3 Discussion of Results

3.3.3.1 \( R_e = R_{ce} = 10 \)

As given in Tables (3.3.2.2), (3.3.2.3) and (3.3.2.4) all attempts to obtain results for the wall vorticity resulted in average percentage errors greater than 100%. The best case was obtained by using the modified boundary condition given by equation (3.2.4.1). However, the least number of iterations was obtained by using the modified inlet correction of Cliffe et al [III.4].
There are several explanations for the large average percentage errors. The first reason comes from the method of calculating the error. According to equation (3.3.1.1) the average percentage error is obtained by dividing by the benchmark values obtained by Cliffe et al [III.4]. The largest errors occur in the separated flow region because the value of the wall vorticity there is very small. Division by a small number leads to large error, especially in the recirculating region.

The second reason for the large average percentage error comes from the method of solution. In Section (2.2) it was noted that the Jacobian J must be greater than zero. In a region of recirculating flow, J can be either greater than zero or less than zero. Hence, the fluid cannot flow along streamlines $\psi = \text{constant}$ in the direction of increasing $\phi(x,y) = x$, as noted in equation (2.3.1). At points in a recirculating region the flow along streamlines $\psi = \text{constant}$ will be alternating between the direction of decreasing $\phi(x,y) = x$ and the direction of increasing $\phi(x,y) = x$.

Although the recirculating region is small, it is still significant enough to cause destabilizing effects on the solution. As seen in Figure (3.3.2.4), the plot of the wall vorticity, the best results are obtained after the recirculating region where errors were as small as 0.1% to 3.0%. As well, as seen in Figure (3.3.2.5), the plot of the worst case wall vorticity obtained by all the participants that reported in [III.3], the values of wall vorticity obtained in this study are quite comparable to results of other investigators.

Numerous attempts were made to improve the results. Clustering the grid according to equations (3.2.5.1) and (3.2.5.2) produced at best oscillating results which would not converge within the given error tolerances. As well, no improvement in the average percentage error was indicated. Attempts to minimize the effects of the recirculating region by shrinking its width, also met with no improvement. For example,
choosing \( R_e = 0.1 \) and \( R_{ee} = 300 \) resulted in converging solutions, but no improvement in the average percentage error - the error was simply spread out more evenly among all the values of \( x \). Also, the limit of Stokes flow (\( R_e = 0; \) infinite viscosity) was attempted for the case under consideration, namely \( R_{ee} = 10 \). In this limit, separation should not occur. The solution did converge, leading to the conclusion that viscous flow with no separation ought to be further investigated by the current method of solution.

Finally, predictions of the separation and reattachment points varied from a low of 13.2\% to a high of 25.2\% error. The separation point was best predicted using the inlet correction of Cliffe et al [III.4] and the reattachment point was best predicted using the boundary condition correction at the inlet.

3.3.3.2 \( R_e = R_{ee} = 100 \)

Converged results for this case could not be obtained. At first it was thought that the problem had to be with the vorticity term in equation (2.4.1.1) because of the larger Reynolds number \( R_e = 100 \). Since the results were oscillating but not converging, it was thought that upwind differencing of the vorticity term in (2.4.1.1) would prevent the oscillations from occurring.

A source of differences among algorithms is in the treatment of the derivatives of the convective terms. In the case where diffusion dominates convection, the use of second order accurate centered differences for the convective derivatives constitutes the best compromise between computational accuracy and economy. In cases where there is convective dominance, the use of centered differences for the convective terms may result in instability or non-physical/oscillatory behaviour. The most widely used way of avoiding this instability is the use of a first order accurate upwind formulation, so that a hybrid central/upwind differencing scheme is employed. Even though this approach
results in unconditional stability, it can introduce artificial or numerical diffusion, an error that may become so dominant as to obscure the effects of the physical diffusivity on the flow. In any event, the oscillations for the case \( R_e = R_{\infty} = 100 \) could not be prevented with this approach.

According to [III.6], there are limitations on the cell Reynolds number, or Peclet number, \( R_e_{cell} \) for non-oscillatory convergent solutions. In [III.6], \( R_e_{cell} \) must be less than or equal to a typical value of 10.

\[
\frac{R_e_{cell}}{\Delta x} \quad \text{constant} \quad \text{and} \quad \frac{R_e_{cell}}{\Delta \psi} \quad \text{constant},
\]

so that taking their ratio, we obtain

\[
\frac{R_e_{cell}}{R_e_{cell}} = \frac{\Delta x}{\Delta \psi} = \text{constant}.
\]

This constant can be looked at as a cell aspect ratio, AR, and should be less than 1. For the case \( R_e = R_{\infty} = 10 \), \( \text{AR} = \frac{0.167}{0.05} = 3.34 \), which is near the value 1 and the method converged. For the case \( R_e = R_{\infty} = 100 \), \( \text{AR} = \frac{1.667}{0.05} = 33.34 \) which is much greater than 1 and the method would not converge. Even though we were restricted by the geometry of the test problem, going to larger values of \( \Delta \psi \) led to solutions that would eventually converge, but the results were not comparable to the benchmark values.

### 3.3.4 Conclusions

Contrary to what was reported by Napolitano and Orlandi in [III.3], this problem was not too easy. In fact, the proposed method of solution in this thesis could not handle the problem, resulting in average percentage errors in excess of 100% for the solutions
that would converge. From these results, it appears that viscous flow problems that contain separated regions of flow cannot be treated, especially at the higher Reynolds numbers, using the formulation in this thesis. However, as we will show in Chapter V, this formulation can be modified in such a way as to allow flow separation. In the next chapter we solve a flow problem which illustrates that the difficulties encountered in this chapter are related to the viscous effects, rather than in handling flow fields with non-zero vorticity.
CHAPTER IV

CIRCULAR CYLINDER IN HYPERBOLIC-COSINE SHEAR FLOW

4.1 INTRODUCTION

Since the problem of recirculation is not fully resolved in Chapter III, a problem is solved where the flow is rotational, but no viscosity is allowed. Here, recirculating flow is eliminated. The problem is hyperbolic-cosine shear flow about a circular cylinder as discussed by Van Dyke [IV.1].

4.2 FLOW EQUATIONS

4.2.1 Differential Equations

If no viscosity is allowed, we are considering flow in the inviscid limit of \( R_e \rightarrow \infty \).

The equation (2.3.7b), namely,

\[
L\{\omega\} = R_e y_\psi \omega_x - y_\psi^2 \omega_\psi = 0
\]

where \( \omega = \omega(x, \psi) \), reduces to \( y_\psi \omega_x = 0 \). Since \( y_\psi = 1/u \neq 0 \), \( \omega_x = 0 \) which implies \( \omega = \omega(\psi) \) only. This condition guarantees that at any \( x \) station (i.e., \( x = \text{constant} \)), the vorticity profile is exactly the same as at the inlet (i.e., \( \omega = \text{constant along each streamline} \)). Thus, once we know \( \omega \) at infinity, we know it throughout the flow field. This is a dynamical condition for steady motion.

Hence, equation (2.3.8c),

\[
L\{y\} - y_\psi^2 \omega_\psi = 0
\]
where \( y = y(x, \psi) \) or,

\[
y^2 y_{xx} - 2y_y y_{x\psi} + (1 + y^2) y_{\psi \psi} - y^2 \omega y_\psi = 0
\]

or in difference form,

\[
\left( A^{(\omega)}_{ij} \delta_{xx} + B^{(\omega)}_{ij} \delta_{x\psi} + C^{(\omega)}_{ij} \delta_{\psi \psi} + E^{(\omega)}_{ij} \delta_\psi \right) y^{(n+1)} = 0
\]

(4.2.1.1)

where the coefficients and subscripts are as previously defined, becomes the only equation to be solved for the unknown \( y = y(x, \psi) \). (Note: We have taken \( \phi = y(x, \psi) \), i.e., \( \alpha = 0, k = n \) in equation (4.2.1.1)). The vorticity \( \omega \) in the coefficient \( E_{ij} \) is known at each \((i, j)\). Again, we can start with some initial approximation of the vorticity \( \omega \) for \( n = 0 \), e.g., \( \omega = 0 \) when \( n = 0 \), just to get the iteration for \( y \) started, afterwards \( \omega = \omega(\psi) \) is known along each streamline.

4.2.2 Circular Cylinder in Hyperbolic-Cosine Shear Flow

The solution of the circular cylinder in hyperbolic-cosine shear flow is developed in Appendix L from the uniform flow problem at infinity to demonstrate the problem’s increasing analytic complexity and which the numerical method utilized in this thesis handles very nicely.

Note that we are not solving an irrotational inviscid problem, but a rotational, inviscid problem because of the upstream boundary condition on vorticity at \( x = -\infty \). This is not a truly physical situation, but is used to simulate an inviscid non-recirculating flow in which vorticity is present.

For flow over a symmetric body, like a circular cylinder, with the flow at infinity symmetric about \( y = 0 \), we only need to consider the upper half \((x, \psi)\) plane. The symmetric body is taken as a section of the streamline \( \psi = 0 \).
4.2.3 Boundary Conditions in the Computational Domain

Given the symmetric body geometry, i.e., \( y = f(x) \), the boundary conditions associated with equation (4.2.1.1) can be expressed as follows (see Diagram (4.2.3.1)).

The values of \( x \) at the leading and trailing edges are denoted by \( x_{LE} \) and \( x_{TE} \), respectively.

Boundary conditions for \( y \): (nondimensionalized):

1. Hyperbolic-cosine shear flow at infinity (equation (L3) from Appendix L):

\[
y = \frac{1}{\epsilon^2} \sinh^{-1}\left(\frac{1}{\epsilon^2}\psi\right) \text{ at } x = \pm \infty \text{ and at } \psi = +\infty
\]

where \( \epsilon = \text{vorticity number (perturbation parameter)} \).

2. Flow symmetry and flow tangency:

\[
y(x,0) = \begin{cases} 0 & \text{for } -\infty < x < x_{LE} \quad x_{TE} < x < \infty \\ f(x) & \text{for } x_{LE} \leq x \leq x_{TE} \end{cases}
\]

Boundary conditions for \( \omega \): (non-dimensionalized).

At all \( x \) stations, \( \omega = \text{constant along a streamline (equation (L4) from Appendix L)} \):

\[
\omega = -\epsilon^2 \sinh(\epsilon^2 y)
\]

for all \( x \).

In discrete form, this is

\[
\omega_{ij} = \omega_{ij} = -\epsilon^2 \sinh\left(\frac{1}{\epsilon^2} y_{ij}\right)
\]
for all j where i = 1 corresponds to x = -∞.

The speed is given by

\[ q^2 = u^2 + v^2 - \frac{1+y_x}{y^2} \]  \hspace{1cm} (4.2.3.1)

We calculate q = q_s, the speed on the surface of the circular cylinder, using equation (4.2.3.1)

\[ q^2 = q_s^2 - \frac{1+f'(x)^2}{y^2} \]  \hspace{1cm} (4.2.3.2)

In discrete form the surface speed is q_{ai} for all i, where j = 1 corresponds to the surface of the circular cylinder.

4.3 RESULTS AND DISCUSSION

4.3.1 Results

Results were obtained for the following:

1. Equal uniform grid spacing in the x and ψ directions, Δx = Δψ = 0.05.
2. The dimensionless vorticity number or perturbation parameter ε = 0.1.
3. The radius of the circular cylinder was r = a = ½.

All results were obtained using an IBM-PC-compatible/Intel 80286 (AT) computer. The numerical results for the speed on the surface of the circular cylinder in hyperbolic-cosine shear flow are presented in the following tables and figures and compared with the second order solution determined by Van Dyke [IV.1] (perturbation or analytic (approximate) solution obtained using equation (L2) in Appendix L).

Table (4.3.1.1) contains a listing of the various parameters used to obtain the results.

Tables (4.3.1.2) and (4.3.1.3) contain the numerical results compared to the perturbation or analytic results obtained from Van Dyke. The error at each point was
compared (the total error was not a consideration in this problem).

Figures (4.3.1.1) and (4.3.1.2) contain a plot of the results tabulated in Tables (4.3.1.2) and (4.3.1.3).

4.3.2 Discussion of Results

As seen either from Tables (4.3.1.2) and (4.3.1.3) or Figures (4.3.1.1) and (4.3.1.2), agreement is best over the main body of the circular cylinder with accuracy decreasing towards the leading and trailing edges. The validity of all these solutions can be questioned near the stagnation points at the leading and trailing edges. As noted in [I.3], the inaccuracy at the leading and trailing edge could be overcome by an extrapolation of the accurate values over the centre of the profile to the stagnation points at the leading and trailing edges. Also, refining the grid near the leading and trailing edges or using a staggered grid spacing would improve the accuracy at these two locations. However, since the main reason for attempting this problem was to validate the method of solution in regions of non-recirculating flow, no further refinements were attempted.

4.3.3 Conclusions

The numerical results are found to be almost exactly the same as the perturbation results except near the leading and trailing edges. The results obtained in [I.3] indicated that the numerical results consistently underestimated the perturbation results. This was not the case for the results obtained here. Making the grid finer or packing lines which would improve the numerical solution in the sense of making it closer to the perturbation results is suggested. However, the perturbation results are themselves only approximate analytic solutions. The numerical results represent the flow accurately in regions not near singularities, i.e., over the top of various profiles, while the perturbation results can
be used to advantage when singularities arise. Hence, we conclude that the numerical results agree very well with the analytical results.

The results of this chapter confirm our assertion that the difficulties encountered in Chapter III are due to the flow separation rather than the flow vorticity. In the next chapter we propose a method which allows treatment of the re-circulating region.
CHAPTER V

STEADY FLOW PAST A BACKWARD FACING STEP

5.1 INTRODUCTION

The flow over a backward facing step (BFS) in a channel or the flow through a channel containing a sudden expansion or corner provides an excellent test case for the accuracy of numerical methods because of the dependence of the reattachment length of the dividing streamline on the Reynolds number $R_e$. Excessive numerical smoothing in favour of stability will result in failure to predict the correct reattachment length.

As noted in [V.1], experimental studies over a BFS yielded two-dimensional laminar flows only at Reynolds numbers $R_e < 400$ and $R_e > 6000$. In the laminar range, the velocity field was close to that of a fully developed channel flow with only a slight deviation from that of parabolic flow.

In between these Reynolds numbers, the transition to turbulent flow was found to be strongly three-dimensional where velocity fluctuations began to increase, while maintaining symmetry to the centre plane of the test section. It was initially believed that the BFS flow, with its simple geometry, would yield a simple flow pattern showing a single separation region attached to the step. Other regions of detached flow were not expected.

Although numerical prediction procedures encounter false diffusion as discussed in Chapter II of this thesis, good agreement between the predicted and measured flow field for Reynolds numbers $R_e < 400$ were obtained in [V.1], demonstrating that truncation errors due to false diffusion can be kept very low. For Reynolds numbers $R_e > 400$, two-dimensional predictions were also obtained, but the results showed multiple
regions of recirculating flow. Some of the differences were explained as being caused by the three-dimensionality of the flow, where additional recirculating flow regions were measured downstream of the primary region of separation caused by the sharp change in the flow direction. For example, an additional recirculating flow region was measured at the upper wall downstream of the expansion, developing late in the laminar range \( R_e = 400 \) and remaining in existence throughout the transition region \( 400 < R_e < 6000 \). This was due to the adverse pressure change or gradient downstream of, and created by, the sudden expansion. It was largely dependent on the expansion ratio of the BFS flow geometry.

For the above reasons, it was decided to try to duplicate the results obtained in [V.2] with a Reynolds number of \( R_e = 50 \). See Diagrams (5.1.1) and (5.1.2) for the physical and computational domains in terms of nondimensional coordinates \( x = \frac{x}{L} \) and \( y = \frac{y}{H} \), i.e., the coordinates have been nondimensionalized in terms of the channel length \( L \) and height \( H \).

5.2 TEST PROBLEM

5.2.1 Problem Specification

Fully developed Poiseuille flow has been specified at the entrance and the exit of the channel. This means, in theory, that equilibrium flow exists at the inlet and exit. Hence, the boundary conditions are such that an imposed fully developed parabolic velocity distribution is prescribed at a section \( x = 0 \) a short distance upstream of the step and also again far downstream at a section \( x = x_{\text{MAX}} \). In reality, the downstream reversion to Poiseuille flow at low \( R_e \) is achieved in an asymptotic manner (as also is the change upstream). This sort of downstream boundary condition becomes invalidated at
high Reynolds number $R_e$ when ultimately turbulence sets in.

As discussed in [V.1], the length measured from the step to the end of the calculation domain should be selected to be equivalent to at least four times the experimentally measured reattachment length of the primary recirculatory region. The boundary condition at that far downstream cross-section can then be taken as that of fully developed parabolic flow, i.e., $\frac{\partial u}{\partial x} = \frac{\partial v}{\partial x} = 0$. This distance has been shown to be sufficient to make the reattachment length independent of the length of the calculation domain. For small Reynolds number, the section length downstream of the step, at the outlet to the channel can be made sufficiently long to permit the flow to redevelop into a fully developed channel flow. For higher Reynolds numbers, no matter how reasonably long the channel, small deviations could be present.

Similar to the discussion by Roache [I.1], the adequacy of the boundary conditions on $y = \gamma(x,\psi)$ and $\omega = \omega(x,\psi)$ in this thesis will depend on the Reynolds number of the problem, the differencing method used, and on the initial conditions. The limit of zero upstream wall boundary layer thickness can be simulated by using a "slip" wall condition. The slip wall condition is that used on the solid wall upper surface or lid. The upstream wall, the base and the centreline will also be solid walls.

The upstream inflow boundary cannot have a unique solution since its characteristics will change depending on the physical flow upstream of the inflow cross-section, and upon the separated flow solution itself. The problem is unclear mathematically. For example, it is not clear that one should completely specify the input, lest the elliptic nature of the equations be restricted. Yet, something must be specified. The upstream inflow boundary is partly determined by specifying a boundary-layer inflow velocity profile shape, and partly develops as part of the solution. By
velocity profile, we will mean values of the x-component of velocity suitably normalized, i.e., on the maximum inlet velocity.

The inflow stream function is determined by integration of the second-order Poiseuille flow boundary-layer solution for \( u \). It should be noted that the boundary-layer equations do not correctly represent the flow at low \( R_e \) and we do not suggest that the input velocity profile used represents an accurate solution of the flat plate flow ahead of the base of the channel. It is merely a convenient one-parameter family of \( u \) velocity profiles by which one can study the effects of upstream velocity profile shape on the separated flow. It qualitatively represents a meaningful flow condition [1.1].

We could neglect the details of farther downstream flow continuation and still obtain realistic answers upstream. However, catastrophic instabilities may be propagated upstream from the outflow boundary and destroy the solution. The aim should be to allow the most free flow adjustment at the downstream continuation surface which still gives a solution. However, the safest method from the viewpoint of stability is to completely specify the outflow conditions (which is the approach taken in this thesis). For example, we will not assume \( \omega_{x,j} = \omega_{y,j} \), which is equivalent to stating that vorticity is merely advected out of the mesh region, assuming no viscous production of vorticity between \((X,j)\) and \((Y,j)\). Hence, at the outlet of the channel, a fully developed velocity profile will be specified.

We will force separation at the sharp corner. As indicated in stream function plots contained in [V.3], the extrapolated separation point moves down from or below the sharp corner off the base. The incompressible numerical results show a regular movement of the separation point down the base as \( R_e \) is decreased, in agreement with the well known incompressible result at \( R_e = 0 \) for Stokes flow over a sphere, in which
no separation occurs. For all the different methods of treating the sharp corner vorticity, separation was indicated below the sharp corner, including one method intended to bias the solution toward separation at the corner and two methods intended to force separation at the sharp corner. At $R_e = 100$, the numerical solution in [V.3] indicated separation occurring at somewhere less than one cell height below the sharp corner. This viscous effect was somewhat exaggerated because of the implicit artificial viscosity effect of the upwind differencing used in [V.3]. The location of separation will thus depend on the cell size, and can only be interpreted qualitatively at $R_e = 100$. On the other hand, at $R_e = 0.1$, the artificial viscosity effect is negligible, so the use of $R_e = 50$ in this thesis seems a reasonable compromise.

Failure to preserve conservation can lead to numerical instabilities. To stabilize the calculations while using methods that do not preserve these properties, artificial viscosity is often introduced, either explicitly or implicitly, by using dissipative finite-difference schemes, especially for high $R_e$ flows. For low $R_e$ flows it is possible that a non-conservative scheme can produce a stable solution without artificial viscosity, since the viscous terms are relatively large anyway and can quickly eliminate the error terms introduced. However, it is interesting to note that even though there are not enough cells between the sharp corner and true separation point to accurately resolve the distance, separation is still indicated experimentally between 2 or 3 cells below the corner. Hence, the phenomenon does not appear to be merely an aberration of the computational mesh [V.3]. Thus, it appears that the Stokes flow limit does give separation below the sharp corner. Also, the backward facing step geometry in [V.3] did not necessarily imply a flow configuration devoid of separation. This phenomenon is not investigated further in this thesis and could be a topic for further study. Instead, we will assume indicted

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separation occurring less than a cell spacing below the corner can be ignored, and the dividing streamline (DSL) can be faired into the assumed corner separation.

At the centre of the recirculatory region, or eye of the separation bubble, there is a strong flow reversal. Because of the conclusions established in Chapter II, this region of the flow field will not be investigated using the mathematical formulation in this thesis.

Finally, at a separation (or reattachment) point in a continuum flow, the vorticity is zero. That is, at the separation and reattachment point, there is a stress free velocity profile. Although the method is not used in this thesis, the reattachment point may be determined by locating the point at which \( \omega = 0 \) along a wall. The present formulation provides a more convenient condition for locating the reattachment point, as will be discussed later.

### 5.2.2 Boundary Conditions in the Physical Domain

Consider steady Poiseuille flow given by the equation

\[
\frac{d^2u}{dy^2} - \frac{1}{\mu} \frac{dp}{dx} = -C, \quad v = 0
\]  

(5.2.2.1)

where \( C \) is a constant.

Integrating (5.2.2.1), we have

\[
u - u(y) = -\frac{C}{2} y^2 + a_m y + b_m
\]  

(5.2.2.2)

where \( a_m \) and \( b_m \) (\( m = 1 \) or \( 2 \)) are arbitrary constants of integration. Equation (5.2.2.2) will be applied at both the inlet (\( m=1 \)) and the outlet (\( m=2 \)).

At the inlet \( x = 0 \) we have \( u = 0 \) at \( y = h \). Hence, (5.2.2.2) gives,
\[-\frac{C_1}{2} h^2 + a_i h + b_i = 0 \quad (5.2.2.3a)\]

where \( C_1 \equiv C \).

Also, \( u = 0 \) at \( y = 1+h, \, x = 0 \), which gives

\[-\frac{C_1}{2} (1+h)^2 + a_i (1+h) + b_i = 0 \quad (5.2.2.3b)\]

Solving (5.2.2.3a) and (5.2.2.3b) for \( a_i \) and \( b_i \) we get

\[ a_i = \frac{C_1}{2} (1+2h) \quad \text{and} \quad b_i = \frac{-C_1}{2} (1+h)h \]

Substituting the above into equation (5.2.2.2) we have

\[ u = u_{\text{tot}}(y) = \frac{-C_1}{2} \left[ y^2 - (1+2h)y + (1+h)h \right] \quad (5.2.2.4a)\]

At the outlet \( x = x_{\text{MAX}} \) we have \( u = 0 \) at \( y = 0 \), which gives using (5.2.2.2),

\[ b_2 = 0 \quad (5.2.2.5a)\]

and \( u = 0 \) at \( y = 1+h \), which gives

\[ -\frac{C_2}{2} (1+h)^2 + a_2 (1+h) = 0 \quad (5.2.2.5b)\]

with \( C \equiv C_2 \).

Solving (5.2.2.5b) for \( a_2 \) we get

\[ a_2 = \frac{C_2}{2} (1+h) \]
Substituting the above into equation (5.2.2.2) we have

\[ u = u_{\text{inlet}}(y) - \frac{C_2}{2} \left[ y^2 - (1+h)y \right] \]  
(5.2.2.6a)

One way to choose \( C_1 \) is to require unit mass flux at the inlet. For example, we take

\[ \psi(y = 1+h) - \psi(y=h) = 1 \text{ at the inlet} \]  
(5.2.2.7)

Using (5.2.2.4a) and the definition \( u = \psi \), we get

\[ \psi_{\text{inlet}}(y) = \frac{y}{h} \int_{h}^{y} u_{\text{inlet}}(y) \, dy \]  
(5.2.2.8a)

\[ = -\frac{C_1}{2} \left[ \frac{y^3}{3} - \frac{1}{2}(1+2h)y^2 + (1+h)hy - \frac{h^3}{3} - \frac{h^2}{2} \right] \]

With \( \psi_{\text{inlet}}(y=h) = 0 \) and choosing \( C_1 \) such that \( \psi_{\text{inlet}}(y=1+h) = 1 \) we can solve (5.2.2.7) and (5.2.2.8a) for \( \frac{C_1}{2} = 6 \).

Then equations (5.2.2.4a) and (5.2.2.8a) become

\[ u = u_{\text{inlet}}(y) = -6[y^2 - (1+2h)y + (1+h)h] \]  
(5.2.2.4b)

and

\[ \psi = \psi_{\text{inlet}}(y) = -2y^3 + 3(1+2h)y^2 - 6(1+h)hy + (3+2h)h^2 \]  
(5.2.2.8b)

Similarly, by the conservation of mass, \( C_2 \) must be chosen to give unit mass flux at the outlet. That is, we take

\[ \psi(y = 1+h) - \psi(y = 0) = 1 \text{ at the exit} \]  
(5.2.2.9)

From (5.2.2.6a), using definition \( u = \psi \), we get
\[ \psi_{\text{outlet}}(y) = \int_0^y u_{\text{outlet}} \, dy \]

\[ = -\frac{C_2}{2} \left[ \frac{y^3}{3} - \frac{(1+h)y^2}{2} \right] \]  \hspace{2cm} (5.2.2.10a)

With \( \psi_{\text{outlet}}(y=0) = 0 \) and choosing \( C_2 \) such that \( \psi_{\text{outlet}}(y = 1+h) = 1 \) we can solve (5.2.2.9) and (5.2.2.10a) for \( \frac{C_2}{2} = \frac{6}{(1+h)^3} \).

Then equations (5.2.2.6a) and (5.2.2.10a) become

\[ u = u_{\text{outlet}}(y) = \frac{-6}{(1+h)^3} \left[ y^2 - (1+h)y \right] \]  \hspace{2cm} (5.2.2.6b)

and

\[ \psi = \psi_{\text{outlet}}(y) = -\frac{1}{(1+h)^3} \left[ 2y^3 - 3(1+h)y^2 \right] \]  \hspace{2cm} (5.2.2.10b)

By definition \( \omega = -\nabla^2 \psi \). But \( \psi = \psi(y) \) only at the inlet \( x = 0 \). Hence,

\[ \omega - \omega_{\text{inlet}}(y) = -\psi_{yy} = -u_y = -6 \left[ 2y - (1+2h) \right] \]  \hspace{2cm} (5.2.2.11)

from either equation (5.2.2.4b) or (5.2.2.8b).

Similarly, \( \psi = \psi(y) \) only at the outlet \( x = x_{\text{MAX}} \).

Hence,

\[ \omega = \omega_{\text{outlet}}(y) = -\psi_{yy} = -u_y = \frac{6}{(1+h)^2} \left[ 2y - (1+h) \right] \]  \hspace{2cm} (5.2.2.12)
from either equation (5.2.2.6b) or (5.2.2.10b).

5.2.3 Boundary Conditions in the Computational Domain

Boundary conditions for \( y \):

- **Inlet boundary** where \( x = 0, 0 \leq \psi \leq \psi_{\text{MAX}} = 1 \)

From equation (5.2.2.8b) \( \psi \) is a cubic function of \( y \) given by

\[
\psi = -2y^3 + 3(1+2h)y^2 - 6h(1+h)y + (3+2h)h^2
\]

This can be solved explicitly for \( y \) as a function of \( \psi \) as per the method in Appendix J to give (for \( h = 0.4 \))

\[
y = \cos(\theta + \frac{4}{3} \pi) + \frac{9}{10}
\]

where

\[
\theta = \frac{1}{3} \arccos (1-2\psi)
\]

- **Outlet boundary** where \( x = x_{\text{MAX}} = 6, 0 \leq \psi \leq \psi_{\text{MAX}} = 1 \)

From equation (5.2.2.10b), \( \psi \) is a cubic function of \( y \) given by

\[
\psi = \frac{1}{(1+h)^3} \left[ 2y^3 - 3(1+h)y^2 \right]
\]

This can be solved explicitly for \( y \) as a function of \( \psi \) as per the method in Appendix J to give (for \( h = 0.4 \))

\[
y = \frac{7}{5} \cos(\theta + \frac{4}{3} \pi) + \frac{7}{10}
\]
where
\[ \theta = \frac{1}{3} \arccos(1 - 2\psi) \]

Alternatively, the downstream continuation problem could be treated in a manner following [V.3] where linear extrapolation for \( y \) is used which approximates
\[ v_x = \left[ \frac{y_x}{y'} \right]_x = 0 \text{ or } y_{xx} = 0 \text{ since } \psi = \psi(y) \text{ (or } y = y(\psi)) \text{ at the outlet } x = x_{\text{MAX}}. \]
This gives the difference equation
\[ y_{ix,j} = 2y_{i+1,j} - y_{i-1,j} \]

- **Upper boundary lid** where \( \psi = \psi_{\text{MAX}} = 1, \ 0 \leq x \leq x_{\text{MAX}} = 6 \)
\[ y = 1 + h = 1.4 \]

- **Lower boundary wall** where \( \psi = 0 (i=1) \)
\[ y = \begin{cases} h = 0.4 \text{ for } 0 \leq x \leq x_{\text{SEP}} = 2 \\ 0 \text{ for } 2 + L_x = x_A \leq x \leq x_{\text{MAX}} = 6 \end{cases} \]

For \( 2 = x_{\text{SEP}} < x < x_A = 2 + L_x \) on the dividing streamline (DSL), see section (5.2.4) for an evaluation of \( y = y_{it}. \)

**Boundary conditions for \( \omega \):**

- **Inlet boundary** where \( x = 0, \ 0 \leq \psi \leq \psi_{\text{MAX}} = 1 \)
From equation (5.2.2.11)
\[ \omega = 6[2y - (1 + 2h)] \]
• Outlet boundary where \( x = x_{\text{MAX}} = 6, \ 0 \leq \psi \leq \psi_{\text{MAX}} = 1 \)

From equation (5.2.2.12)

\[
\omega - \frac{6}{(1+h)^2} [2y-(1+h)]
\]

Alternatively, in a manner following [V.3] we could impose zero streamwise gradient for \( \omega \), namely \( \omega_x = 0 \), which has finite difference approximation

\[
\omega_{ix,j} = \omega_{i,j}
\]

• Upper boundary lid where \( \psi = \psi_{\text{MAX}} = 1 \) (\( i = \text{IX}, \ 0 \leq x \leq x_{\text{MAX}} = 6 \))

In [V.3] the boundary condition at the upper boundary was reported as disappointing, resulting in destabilizing solutions. It was not possible to model the backstep with no upper boundary or in the free flight case (inflow through the mesh of the lid). The most nearly free condition in [V.3] was to use an impermeable slip wall at the lid.

Since \( v = 0 \) on \( \psi = \psi_{\text{MAX}} \), the following condition for \( \omega \) at the lid was used in [V.3], namely \( \omega_{i,ix} = 0 \), which implies \( \frac{\partial u}{\partial y} = 0 \) since

\[
\omega = -\frac{\partial u}{\partial y} + \frac{\partial v}{\partial x} - \frac{\partial u}{\partial y} = 0 \quad \text{(since } v = 0 \text{ for all } x \text{ along the lid)}.
\]

Because the wall in our case is assumed not to be impermeable, \( \omega_{i,jx} = \omega_{i,j1} \) is a less restrictive condition, which approximately implies \( \frac{\partial^2 u}{\partial y^2} = 0 \) at the lid since

\[
\frac{\partial \omega}{\partial y} - \frac{\partial^2 u}{\partial y^2} + \frac{\partial^2 v}{\partial x \partial y} - \frac{\partial^2 u}{\partial y^2} = 0 \quad \text{(since } v = v(y) \text{ only)}.
\]
Lower boundary wall where \( \psi = 0 \) (\( i = 1 \))

From equation (2.4.5.7g)

\[
\omega_{i1} = -\frac{1}{12\Delta\psi} \left( 2q_i^2 - 9q_3^2 + 18q_2^2 \right)
\]

for \( 0 \leq x < x_{\text{SEP}} = 2 \) and \( 2 + L_x = x_A < x \leq x_{\text{MAX}} = 6 \)

\( \omega = 0 \) for \( x = x_{\text{SEP}} = 2 \) (modified to allow for the singularity) and \( x = x_A = 2 + L_x \)

For \( 2 = x_{\text{SEP}} < x < x_A = 2 + L_x \) on the DSL we use equation (2.4.5.6a), namely

\[
\omega = v_x - \frac{1}{2} q_2 \frac{\psi^2}{\Delta\psi}
\]

which is approximated as follows:

\[
\omega_{i1} \approx \frac{V_{i+1,1} - V_{i-1,1}}{2\Delta x} - \frac{1}{2} \frac{q_{i2}^2 - q_{i1}^2}{\Delta\psi}
\]

where

\[
v_{k1} \approx \frac{y_{k+1,1} - y_{k-1,1}}{2\Delta x} \cdot \frac{\Delta\psi}{y_k - y_{k1}} \quad \text{for} \quad k-i-1 \quad \text{and} \quad i+1
\]

From equation (2.4.3.1), the speed is approximated using

\[
q_{ii}^2 \approx 1 + \left( \frac{y_{i+1,1} - y_{i-1,1}}{2\Delta x} \right)^2 \left( \frac{y_{i2} - y_{i1}}{\Delta\psi} \right)^2
\]

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and

$$q_{i2}^2 \approx 1 + \left( \frac{y_{i+1,2} - y_{i-1,2}}{2\Delta x} \right)^2 \left( \frac{y_{i3} - y_{i1}}{2\Delta \psi} \right)^2$$

Alternatively, from equation (2.4.5.7g), with $q_{ni}^2 \neq 0$, the second term in $\omega$ could be approximated by

$$\frac{1}{24\Delta \psi} \left( 2q_{i3}^2 - 9q_{i3}^2 + 18q_{i1}^2 - 11q_{ni}^2 \right)$$

As discussed in [V.4], it is worth noting here that as $R_e$ decreases ($\leq 100$) the reattachment point moves forward towards the base. Hence, the vorticity contour lines have a distinct similarity to streamlines indicating less dominance of advective transport. For example, the plots of various DSLs appear similar (they are virtually identical) to the dividing vorticity contours plotted in [V.3] for similar Reynolds numbers. Hence, for ease of computation we could take $\omega = 0$ on the DSL without creating any significant error for the Reynolds number under consideration.

5.2.4 Expression for $y = y_i$ on the Lower Boundary

From the conservation of mass, considering the conduction of constant mass flux applied to the stream tube adjacent to the lower boundary, we obtain

$$\Delta \psi = \int_{y_a}^{y_b} u(x_i, y) \, dy \quad (5.2.4.1)$$
where we take \( y_2 > y_1 \).

Using the Trapezoidal Rule to approximate the integral in (5.2.4.1), we obtain

\[
\Delta \psi \approx \frac{\Delta y}{2} (u_1 + u_2)
\]

Hence, approximately, we have

\[
\Delta \psi = \left[ \frac{y_2 - y_1}{2} \right] \left[ \frac{1}{y_1} \left| \frac{1}{u_1} \left| + \frac{1}{y_2} \left| \frac{1}{u_2} \right| \right. \right] \quad (5.2.4.2)
\]

where we have used \( u = \frac{1}{y_\psi} \) from (2.3.5).

Approximating \( y_\psi \big|_{\text{i/ii}} \) using a 2-point forward difference, \( (\delta_\psi y)_{\text{i/ii}} = \frac{y_2 - y_1}{\Delta \psi} \), and

\( y_\psi \big|_{\text{i/ii}} \) using a 2-point central difference, \( (\delta_\psi y)_{\text{i/ii}} = \frac{y_2 - y_1}{2 \Delta \psi} \), and simplifying, we obtain from equation (5.2.4.2)

\[
y_1 = 2y_2 - y_3 \quad (5.2.4.3)
\]

It should be noted that equation (5.2.4.3) is the 3-point central difference formula for \( (\delta y_\psi)_{\text{i/ii}} = 0 \) and is valid on the lower boundary \( \psi = 0 \) (j = 1) for all \( x \), including the interval \( 2 = x_{\text{sep}} < x < x_A = 2 + L_x \). It is, therefore, an equation for the dividing or separation streamline in the physical plane.

Having obtained \( y_1 \), call these values \( y_{\text{ii}}^{(n+1)} \). A modified or relaxed value \( y_{\text{ii}}^{(n+1)} \) can be obtained by using the relaxation factor \( \beta \), with the following expression:

\[
y_{\text{ii}}^{(n+1)} = (1 - \beta) y_{\text{ii}}^{(n)} + \beta y_{\text{ii}}^{(n+1)} \quad (5.2.4.4)
\]

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The reattachment point, where \( x = x_A = 2 + L_x \), is determined from the value of \( i \) at which \( y_i^{(n+1)} = 0 \). The reattachment length is \( L_x = x_A - 2 \). Note that \( y_i^{(n+1)} = 0.4 \) at \( x = x_{sep} - 2 \) and the initial guess for \( y_i^{(0)} (n=0) \), in order to begin the iteration procedure, was chosen to lie on an ellipse (see Appendix M).

5.2.5 **Modification of Equation (5.2.4.3) for \( y = y_i \) on the Lower Boundary Dividing Streamline**

Equation (5.2.4.3), namely,

\[ y_i = 2y_2 - y_3 \]

expresses the conservation of mass, approximately. Hence, it should hold for all \( i \), in particular, at the inlet.

Since, at the inlet \( \psi = \psi(y) \) only, we have

\[
\Delta \psi = \frac{dy}{dy} \Delta y = u \Delta y \quad (5.2.5.1a)
\]

In the computational domain \( \Delta \psi = \text{constant} \) as \( j \) increases from 1 to \( J1 \), but \( \Delta y \) is not constant. Define

\[
\Delta y_j = y_{j+1} - y_j \quad (5.2.5.2a)
\]

where we have dropped the \( i \) index for convenience.

Rearranging equation (5.2.5.1a) we have

\[
y_{j+1} = y_j + \Delta y_j, \quad j=1,2,\ldots,J1 \quad (5.2.5.2b)
\]
For variable $\Delta y_j$, we can rearrange equation (5.2.5.1a) as

$$\Delta y_j = \frac{\Delta \psi}{u_j}$$  \hspace{1cm} (5.2.5.1b)

Now, as $y_j$ increases (i.e., as $j$ increases), we know $u$ increases in the lower half of the channel. For example, from equation (5.2.2.11)

$$\frac{du}{dy} = -6[2y-(1+2h)] = 6-12(y-h) > 0$$

for $|y-h| < \frac{1}{2}$, i.e., $h \leq y < h+\frac{1}{2}$.

Thus, as $y_j$ increases, $u_j$ increases and hence, $\Delta y_j$ decreases from equation (5.2.5.1b). Therefore, up to the centre of the channel $j = j_c$, we have

$$\Delta y_1 > \Delta y_2 > ... > \Delta y_{j_c}$$  \hspace{1cm} (5.2.5.3)

Now consider equation (5.2.4.3) at the inlet (dropping the $i$ index), namely,

$$y_1 = 2y_2 - y_3$$  \hspace{1cm} (5.2.5.4)

From equation (5.2.5.2b) we have

$$y_1 = h$$
$$y_2 = y_1 + \Delta y_1$$
$$= h + \Delta y_1$$  \hspace{1cm} (5.2.5.5)
$$y_3 = y_2 + \Delta y_2$$
$$= h + \Delta y_1 + \Delta y_2$$
Substituting (5.2.5.5) into (5.2.5.4) we have

\[ h = 2(h + \Delta y_1) - (h + \Delta y_1 + \Delta y_2) \]

which leads to

\[ \Delta y_1 = \Delta y_2 \]

in contradiction to inequalities (5.2.5.3).

The same kind of analysis can be applied at the separation point \( x = x_{\text{sep}} = 2 \).

We would still expect that

\[ \Delta y_1 > \Delta y_2 > \ldots > \Delta y_{s} \text{ at } x = x_{\text{sep}} = 2 \]

Then equation (5.2.5.4), namely

\[
\begin{align*}
y_1 &= 2y_2 - y_3 \\
&= 2(h + \Delta y_1) - (h + \Delta y_1 + \Delta y_2) \\
&= h + \Delta y_1 - \Delta y_2 \\
&> h \text{ since } \Delta y_1 > \Delta y_2
\end{align*}
\]

This means that equation (5.2.4.3), or as modified in (5.2.5.4), predicts that the dividing or separation streamline rises rather than falls just after the separation point \( x = x_{\text{sep}} = 2 \).

Obviously, this is not physically realistic. The problem lies in the fact that constant \( \Delta \psi \)'s correspond to a widening of the streamtubes near the walls, where \( u \) is small, and a narrowing near the mid-stream. This is precisely the opposite of what we want to occur.

To alleviate this problem, define a new variable \( \eta \) by

\[ \eta = \eta(\psi) \quad (5.2.5.6) \]
such that

\[ \psi = \psi_{\text{int}}(\eta) = -2\eta^3 + 3(1+2h)\eta^2 - 6(1+h)h\eta + (3+2h)h^2 \quad (5.2.5.7) \]

Equation (5.2.5.7) is equation (5.2.2.8b) with y replaced by \( \eta \).

The new variable \( \eta \) will be used to replace \( \psi \), that is, we perform the following change of independent variables:

<table>
<thead>
<tr>
<th>(x,y)</th>
<th>( \rightarrow )</th>
<th>(x,\psi)</th>
<th>( \rightarrow )</th>
<th>(x,\eta)</th>
</tr>
</thead>
<tbody>
<tr>
<td>physical domain</td>
<td>intermediate computational domain</td>
<td>final computational domain</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

The overall effect of this sequence of transformations is that we will have uniform spacing \( \Delta x \) and uniform spacing \( \Delta y \) only at inlet in the physical domain, uniform spacing \( \Delta x \) but variable spacing \( \Delta \psi \) in the intermediate computational domain, and uniform spacing \( \Delta x \) and uniform spacing \( \Delta \eta \), in the final computational domain. That is, the effect of uniform or constant spacing \( \Delta \eta \) in the \((x,\eta)\) plane is to give variable spacing \( \Delta \psi \) in the \((x,\psi)\) plane, whereas previously the spacing in the \((x,\psi)\) plane was uniform or constant. Hence, we have effectively packed the grid near the walls.

On the lower boundary \((j=1)\) in the \((x,\eta)\) plane we now have, for all \( i \),

\[ \eta_1 = h \]

and,

\[ \eta_2 = \eta_1 + \Delta \eta = h + \Delta \eta \quad (5.2.5.8a) \]

where \( \Delta \eta = \eta_2 - \eta_1 \).

On the lower boundary \((j=1)\) in the \((x,\psi)\) plane we have, for all \( i \), \( \psi_1 = 0 \), and \( \Delta \psi = \psi_2 - \psi_1 \).
so that $\psi_2 = \psi_1 + \Delta \psi = \Delta \psi \quad (5.2.5.8b)$

Again, from the conservation of mass, considering the condition of constant mass flux applied to the stream tube adjacent to the lower boundary, we obtain

$$\Delta \psi = \int_{y_a}^{y_b} u \, dy$$
$$= \psi_2, \quad \text{from (5.2.5.8b)}$$
$$= \psi(\eta_2)$$
$$= \psi(h + \Delta \eta), \quad \text{from (5.2.5.8a)}$$
$$= (3 - 2\Delta \eta) \Delta \eta^2, \quad \text{using equation (5.2.5.7)}$$

Using $u = \frac{1}{y\psi}$ we have

$$y\psi = \frac{\partial y}{\partial \psi} - \frac{d\eta}{d\psi} \frac{\partial y}{\partial \eta} - \frac{y}{y\psi} \frac{d\psi}{d\eta} - \frac{y}{\psi'(\eta)}$$

Hence, combining the above results,

$$\Delta \psi = \int_{y_a}^{y_b} u \, dy - \int_{y_a}^{y_b} \frac{1}{y\psi} \, dy - \int_{y_a}^{y_b} \frac{\psi'(\eta)}{y\psi} \, dy = (3 - 2\Delta \eta) \Delta \eta^2$$

Using the trapezoidal rule to approximate the above integral, we obtain
\[
\left( \frac{y_{i2} - y_{i1}}{2} \right) \left[ \frac{\psi(\eta)}{y_\eta} \bigg|_{i_1} + \frac{\psi(\eta)}{y_\eta} \bigg|_{i_2} \right] = (3 - 2\Delta\eta)\Delta\eta^2 \quad (5.2.5.9)
\]

From equation (5.2.5.7)

\[
\psi' = \psi(\eta) = -6\left[\eta^2-(1+2h)\eta+(1+h)h\right] \quad (5.2.5.10)
\]

so that \(\psi'(\eta)_{i_1} = \psi'(\eta)_{i_2} = \psi'(h) = 0\).

Therefore, equation (5.2.5.9) reduces to

\[
\left( \frac{y_{i2} - y_{i1}}{2} \right) \frac{\psi'(\eta)}{y_\eta} \bigg|_{i_2} = (3 - 2\Delta\eta)\Delta\eta^2 \quad (5.2.5.11)
\]

Now \(\psi'(\eta)_{i_2} = \psi'(\eta)_{i_2} = \psi'(h+\Delta\eta) = 6(1-\Delta\eta)\Delta\eta, \quad \text{from (5.2.5.10)}\)

Approximating \(y_{i2}\) using a central difference (i.e., \(\left(\delta_y_{i2} - \frac{y_{i2} - y_{i1}}{2\Delta\eta}\right)\) and using the above expression for \(\psi'(\eta)_{i_2}\), we obtain from (5.2.5.11)

\[
y_{i1} = \frac{6(1-\Delta\eta)y_{i2} - (3 - 2\Delta\eta)y_{i3}}{3 - 4\Delta\eta} \quad (5.2.5.12)
\]

Note that if \(\Delta\eta = 0\), equation (5.2.5.12) reduces to equation (5.2.4.3). Numerically, it is possible to find values of \(\Delta\eta\) such that equation (5.2.5.12) predicts a decrease in \(y = y_{i1}\) for \(x \geq x_{\text{sep}} = 2\), i.e., \(y = y_{i1} < h\) for \(x \geq x_{\text{sep}} = 2\). Hence, the dividing or
separation streamline falls rather than rises just after the separation point.

5.3 RESULTS AND DISCUSSION

5.3.1 Preliminaries

For the geometry and data depicted in Diagram (5.1.1), a stream function-vorticity method with equally spaced mesh sizes (=0.1) in both coordinate directions was used in [V.2]. After 328 iterations, the length of the recirculatory region was determined to be \( L_x = 0.765 \). Other results range from a low of \( L_x = 0.580 \) to a high of \( L_x = 0.770 \).

As noted in [V.4], a Reynolds number based on the expansion ratio \( E = \frac{h}{H} = \frac{h}{1+h} \), or step height \( h \), as a single parameter that defines the reattachment length in a laminar two-dimensional flow, is unlikely. The reattachment length in laminar two-dimensional BFS flows is probably not a function of a single variable, but more likely a function of several variables, including the expansion ratio and the inlet section Reynolds number \( R_c \) (based on the maximum inlet velocity and twice the inlet channel height or twice the hydraulic radius of the inlet or small channel \( 2(H-h) \)).

The relevant limits for the examination of the planar laminar BFS flow field are \( E \rightarrow 0 \) and \( R_c \rightarrow 0 \), for which \( L_x \rightarrow 0 \). Experimental evidence indicates that the flow field becomes three-dimensional and turbulent for sufficiently large values of \( E \) and \( R_c \) [V.4]. The best parametric fit for the range of expansion ratios and Reynolds numbers covered by data in [V.4] is given by a correlation of the form

\[
\frac{L_x}{2(H-h)} = C_1R_c^{C_2}(e^{C_3E}-1)
\]

(5.3.1)
where \( C_1 = 0.004, C_2 = 0.75 \) and \( C_3 = 4.75 \). Again, for the geometry depicted in Diagram (5.1.1), equation (5.3.1) with \( E = \frac{h}{H} = \frac{0.4}{1.4} \) and \( R_e = 50 \) gives \( L_x = 0.434 \).

The above should give some indication in the variability of the results obtained for the length of the primary recirculating region.

### 5.3.2 Results

Results were obtained for equal-sized grid spacing in the \( x \) and \( \eta \) directions, i.e., \( \Delta x = \Delta \eta = 0.1 \). All results were obtained using an IBM PC-compatible/Intel 80286 (AT) computer.

The numerical results for the solution of the flow over a BFS in a channel or the flow through a channel containing a sudden expansion or corner are presented in the following table and figure.

Table (5.3.2.1) contains a listing of the various parameters used to obtain the results.

Figure (5.3.2.1) contains a plot of the streamlines \( y_{ij} \), for \( j = 1 \), the lower boundary, (the upstream wall surface, dividing streamline and downstream wall surface) and \( j = 2 \) (first streamline) to \( j = 10 \) (upper surface) for \( i = 1,2,\ldots,60 \).

Figure (5.3.2.2) contains a plot of the vorticity distribution \( \omega_{ij} \) along the lower boundary (upstream and downstream walls and dividing streamline, i.e., \( j = 1 \)), and the upper wall (\( j = 10 \)), for \( i = 1,2,\ldots,60 \).

Using the original boundary conditions described in section 5.2.3 rather than the alternative ones from (V.3), after the 11th iteration, the dividing or separation streamline (DSL) crossed the \( x \) axis for the first time. After 363 iterations, the value of \( L_x \) fell between 0.75 to 0.80 (\( x = 2.75 \) to 2.80). Using the mid-value for \( L_x \) (0.755), this gave
an error of 1.3% when compared to the results obtained in [V.2]. Similar (but not as accurate) results were obtained using various combinations of the alternative boundary conditions.

5.3.3 Discussion of Results

Difficulties were encountered in achieving global convergence. Although it did not seem to matter which combination of boundary conditions were used, as described in section (5.2.3), the vortex length or reattachment length $L_x$ simply oscillated back and forth about the nominal value reported in the previous section, regardless of the number of iterations. As noted in [V.5] and [V.6] the use of smooth boundaries enabled the authors to remove the vorticity singularities at sharp corners at which the boundary slope is discontinuous. At the corners the vorticity is infinite. This was done by using either a boundary curve with a continuous slope or replacing the corner by a blunt stagnation point. The model with smooth boundaries gave the authors in [V.5] and [V.6] a more realistic representation of the physical flow being modelled. In our case, following this lead, the singularity in vorticity at the step $x=2$ was removed using a boundary curve with a continuous slope, and global convergence was achieved. Overall results for $y_{ii}$ changed very little, probably because the stream function is not singular anywhere so an alternative method is only needed for deriving the vorticity at points where the difference equations employ values of $\omega$ at the singular points themselves as noted in [V.5].

As noted in [V.4], at a specified expansion ratio $E$, $L_x$ grows in a moderately nonlinear manner as $R_e$ increases. That is, in the laminar flow regime, the length of the flow development or reattachment downstream of the step increases with increasing Reynolds number; however, the increase is not linear. At a given or constant $R_e$, a monotonic increase (nearly exponentially) in $L_x$ occurs with increasing $E$, because an
increase in E causes an increase in the maximum velocity at the inflow or channel inlet. Both of the above effects are observed in our case when the parameters E and \( R_e \) are varied.

As noted in [V.4], the DSL appeared concave upwards near reattachment. This result is not observed in our case, where the results for \( y_i \) simply decrease steadily for increasing \( x \).

Reattachment occurs between one and two base heights downstream of the corner for the input velocity problem (second order) [V.3]. Although not attempted in this thesis, if the order of the input profile increases, reattachment occurs further downstream of the corner, i.e., for a fourth order profile, reattachment should occur between seven and eight base heights downstream.

Finally, it should be noted that following Chapter IV, when the formulas for \( \omega_i \) were removed, global convergence also resulted, indicating that there is possibly some instability created in the use of the formulas reported in equations (2.4.5.7a) through (2.4.5.8i). This problem with the approximations used for boundary vorticity values has been noted by several authors, for example, see [II.1].

5.3.4 Conclusions

Once the correct and accurate formulas are obtained for locating the position of the dividing streamline which bounds the recirculation zone, and the problem of singularity in vorticity at the corner of the BFS is removed, the numerical results are found to be almost exactly the same as the results presented in [V.2]. Problems with discontinuities in the flow geometry resulting in singularities in one of the flow parameters are to be avoided unless special care is taken to handle the difficulties that will arise.
CHAPTER VI

CONCLUSIONS

The work in this thesis was undertaken to investigate the feasibility, practicality and advisability of applying stream function coordinate methods in viscous flow problems. The natural curvilinear coordinate system \((\phi, \psi)\) introduced by Martin [1.2] and the use of von Mises coordinates \((x, \psi)\) introduced by Barron [1.4] has been extended to the consideration of two-dimensional, steady, incompressible, laminar viscous flows.

In Chapter III, difficulties associated with a weakly separated region in the flow field are discussed. It has been shown, by considering the well-documented test problem of flow in a smoothly expanding channel, that the stream function coordinate method in its conventional form cannot accurately predict viscous flows if the flow has separated. The question arises, however, as to whether the difficulties are associated with the numerics or the physics. In particular, the failure to achieve an accurate converged solution could be due to the sensitivity of the vorticity transport equation solution and/or the choice of numerical approximation for the wall vorticity (numerics), or the presence of viscosity which leads to no-slip conditions on the solid walls (physics). Nevertheless, the conclusion can be drawn that flow separation leading to the development of a recirculating region must be handled carefully, or completely avoided if possible. It was further observed that the ratio of the grid spacing \(\frac{\Delta x}{\Delta \psi}\) should not be too large, preferably as near as possible to one as the calculations will allow.

The study of an inviscid shear flow over an obstacle in Chapter IV indicates that the primary difficulty in applying von Mises coordinates is related to the viscosity and the recirculating flow region rather than the presence of vorticity in the flow.
In Chapter V, problems in which the flow configuration results in singularities in one of the flow parameters are discussed. The classical problem of flow over a backward facing step is considered. The stream function coordinate method is modified to account for flow separation and reattachment. A novel scheme is devised, based on conservation of mass flux, to accurately predict the location of the dividing streamline and its point of reattachment to the wall.

This thesis successfully demonstrates the application of the theory as developed in Chapter II and represented by equations (2.3.8c) and (2.3.7b). Stream function coordinate methods can be used for both attached and separated flows provided appropriate attention is paid to modelling the separation and recirculating region. The method cannot be used to study the recirculating flow itself. However, it does allow one to accurately identify the region of recirculation which can then be solved by conventional methods. The importance of the stream function coordinate method for viscous flows is that it allows efficient grid usage since the main part of the flow field can be predicted by this method. Numerical grid generation is only required in the relatively small recirculating regions which can be tightly gridded and accurately resolved using conventional formulations and methods.
REFERENCES


Figure (3.3.2.1) Streamlines With No Inlet Correction ($R_e = R_{\infty} = 10$)
Figure (3.3.2.2) Streamlines With Inlet Correction (R_e = R_{\infty} = 10)
Figure (3.3.2.3) Streamlines With Boundary Condition Correction \((R_e = R_\infty = 10)\)
Figure (3.3.2.4) Wall Vorticity ($R_e = R_{\infty} = 10$)
Figure (3.3.2.5) Worst Case Wall Vorticity Values in [III.3] vs. Results in Figure (3.3.2.4) ($R_e = R_{ec} = 10$)
Figure 4.3.1.1
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Figure (4.3.1.2)  Speed on Surface of Circular Cylinder: Analytic Solution Using Equation (4.2.3.2) (Van Dyke’s Perturbation Solution) vs. Numerical Solution
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Figure (5.3.2.2) Vorticity Distribution for BFS
Table (3.3.2.1) Parameters for $R_c = R_{\infty} = 10$

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**Average Percentage Error**: 124.5%

**Number of Iterations**: 381
Table (3.3.2.3) Vorticity with Inlet Correction
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<th>Vorticity</th>
<th>Relative Error x 100%</th>
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AVERAGE PERCENTAGE ERROR 138.5%

NUMBER OF ITERATIONS 365
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AVERAGE PERCENTAGE ERROR 123.3%

NUMBER OF ITERATIONS 375
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Table (4.3.1.1) Parameters for $R_c = \infty$

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Number of Iterations 351
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Number of Iterations: 341
Table (5.3.2.1) Parameters for $R_\gamma = 50$

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<tr>
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<td>6.0</td>
</tr>
<tr>
<td>$\psi_{\min}$</td>
<td>0.00</td>
</tr>
<tr>
<td>$\psi_{\max}$</td>
<td>1.00</td>
</tr>
<tr>
<td>$\delta$</td>
<td>0.40</td>
</tr>
<tr>
<td>$\beta_\gamma$</td>
<td>1.80</td>
</tr>
<tr>
<td>$\beta_u$</td>
<td>0.20</td>
</tr>
<tr>
<td>$\varepsilon_\gamma$</td>
<td>$1 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\varepsilon_u$</td>
<td>$1 \times 10^{-2}$</td>
</tr>
<tr>
<td>$\Delta \eta$</td>
<td>0.17-0.18</td>
</tr>
</tbody>
</table>
DIAGRAMS
Diagram (2.1.1) \((\phi, \psi)\) Coordinate System
Diagram (3.1.1) Physical Domain

- Symmetry line: $y = y_s(x) = 1$
- Outlet: $(\frac{1}{3}R_{cc}, 1)$
- Inlet: $(0, 0)$
- Wall: $y = y_l(x)$
- Boundary: $(\frac{1}{3}R_{cc}, y_l(\frac{1}{3}R_{cc}))$
Diagram (3.1.2) Computational Domain

\[ y = y_e(x) = 1 \]

\[ (0, 1) \quad \quad \quad \quad \text{Inlet} \]

\[ y = y_i(x) \]

\[ (0, 0) \quad \quad \quad \quad \quad \text{Outlet} \]

\[ \left( \frac{1}{3} R_{cc}, 1 \right) \]

\[ \left( \frac{1}{3} R_{cc}, 0 \right) \]
Diagram (L1) Slight Shear Flow Past a Circular Cylinder
Diagram (4.2.3.1) Computational Domain

\[ L(y) - y_\psi \omega y_\phi = 0 \]

For a circular cylinder of radius \( r = a = \frac{1}{2} \), \( y = f(x) = \sqrt{0.25 - x^2} \)
Diagram (5.1.1) Physical Domain

\[ R = \text{recirculating region} \]

Inlet: \[ u = \text{equation (5.2.2.4b)} \quad v = 0 \]

Length of channel = 6

Height of channel = \( H = 1 + h = 1.4 \)

Outlet: \[ u = \text{equation (5.2.2.6b)} \quad v = 0 \]

Backward Facing Step (BFS) located at \( x = 2 \)

Height of BFS \( h = 0.4 \)

Length of \( R \) vortex length or reattachment length = \( L_x \)

Reynolds number \( R_e = 50 \)

Expansion ratio \[ E = \frac{h}{H} = \frac{h}{1+h} \]
Diagram (5.1.2) Computational Domain

\[ \psi_{\text{MAX}} = 1 \]

\[ y = 1 + h \]

\[ y = h \quad \text{DSL} \quad y = 0 \]

Inlet

\[ y = y(\psi) = \text{equation (5.2.2.8b)} \]

Outlet

\[ y = y(\psi) = \text{equation (5.2.2.10b)} \]
APPENDICES
APPENDIX A

Presentation of equations (2.2.2c) and (2.2.3b) using equation (2.2.6) in an alternate form

From equations (2.2.2c)

\[ h_{xy} - (v \omega)_y = -\frac{1}{R_e} \omega_{yy} \]
\[ h_{yx} + (u \omega)_x = \frac{1}{R_e} \omega_{xx} \]

Subtracting, we get

\[ \frac{1}{R_e} \nabla^2 \omega = (u \omega)_x + (v \omega)_y \]
\[ = (u_x + v_y) \omega + u \omega_x + v \omega_y \]
\[ = u \omega_x + v \omega_y \]

using (2.2.1b)

Hence, using (2.2.6), \( \nabla^2 \omega + R_e (\psi_x \omega_y - \psi_y \omega_x) = 0 \).

From equation (2.2.3b), using (2.2.6) again, we get

\[ \omega = -\nabla^2 \psi \]  \hspace{1cm} (A1)
APPENDIX B

Derivation of equations (2.2.2e) from (2.2.2d)

Equations (2.2.2d) are

\[ Gh_\phi - F(h_\phi + \omega) = -\frac{1}{R_e} J_\omega \]  
(B1)

\[-Fh_\phi + E(h_\phi + \omega) = -\frac{1}{R_e} J_\omega \]  
(B2)

Dividing (B1) by G, (B2) by F, and adding gives

\[(h_\phi + \omega) \left( \frac{F}{G} - \frac{E}{F} \right) - \frac{1}{R_e} \frac{J}{G} \omega_\phi + \frac{1}{R_e} \frac{J}{F} \omega_\phi = 0 \]  
(B3)

But \( \frac{F}{G} - \frac{E}{F} = -\frac{J^2}{FG} \). Dividing (B3) by this quantity, and noting that

\[ \left( \frac{-J}{G} \right) \left[ -\frac{FG}{J^2} \right] = \frac{F}{J} \]

we get

\[ h_\phi = -\omega + \frac{1}{R_e} \frac{G}{J} \omega_\phi - \frac{1}{R_e} \frac{F}{J} \omega_\phi \]  
(B4)

From (B1) and (B4)

\[ h_\phi = \frac{1}{R_e} \frac{F}{J} \omega_\phi - \frac{1}{R_e} \left[ \frac{F^2}{GJ} + \frac{J}{G} \right] \omega_\phi \]

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Using
\[ \frac{F^2}{GJ} + \frac{J}{G} - \frac{F^2 + EG - F^2}{GJ} = \frac{E}{J} \]
yields
\[ h_\phi = \frac{1}{R_e \frac{J}{J} \omega_\phi} - \frac{1}{R_e \frac{J}{J} \omega_\phi} \]
APPENDIX C

Demonstration that Gauss’ equation (2.2.7) is automatically satisfied

We have, from (2.2.7),

\[
\left[ \frac{J}{E} \Gamma_{11}^2 \right]_x - \left[ \frac{J}{E} \Gamma_{12}^2 \right]_x = 0 \quad \text{(C1)}
\]

By definition,

\[
\Gamma_{11}^2 = \frac{-FE_x + 2EF_x - EE_y}{2J^2} = \frac{-y_x y_x^2 y_x y_{xx} + 2(1 + y_x^2)(y_x y_{xx} + y_x y_{xx}) - (1 + y_x^2)2y_x y_{xx}}{2y_x^2} = \frac{y_{xx}}{y_x}
\]

if we assume that \( y_x = y_{xx} \).

Hence,

\[
\frac{J}{E} \Gamma_{11}^2 = \frac{y_x}{1 + y_x^2} \quad \frac{y_{xx}}{1 + y_x^2}
\]

Also, by definition,

\[
\Gamma_{12}^2 = \frac{F \sigma_x - FE_y}{2J^2} = \frac{(1 + y_x^2)2y_x y_{xx} - y_x y_x^2 y_x y_{xx}}{2y_x^2} = \frac{y_{xx}}{y_x}
\]

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Hence,

\[ \frac{J}{E} \Gamma^{2}_{11} \phi \frac{y_{\phi}}{1+y_x^2} y_{\phi x} \frac{y_{\phi x}}{1+y_x^2} \]

Substituting into the left-hand side of equation (C1) gives

\[ \left[ \frac{J}{E} \Gamma^{2}_{11} \right] \phi \frac{y_{\phi}}{1+y_x^2} y_{\phi x} \frac{y_{\phi x}}{1+y_x^2} \left[ \frac{y_{xx}}{1+y_x^2} \right] \phi \left[ \frac{y_{\phi x}}{1+y_x^2} \right] x \]

\[ - \frac{(1+y_x^2)y_{xx\phi} - 2y_xy_{xx}y_{x\phi} - (1+y_x^2)y_{\phi xx} + 2y_xy_{xx}y_{x\phi}}{(1+y_x^2)^2} \]

\[ - 0 \]

if we assume that \( y_{x\phi} = y_{\phi x} \) and \( y_{xx\phi} = y_{\phi xx} \) for all \( (x, \psi) \).
APPENDIX D

Derivation of equation (2.3.7b) from equations (2.3.7a)

From the first equation of (2.3.7a), differentiating with respect to \( \psi \),

\[
R_c h_{x\psi} = y_x \omega_x + y_{xx} \omega_x - \frac{(1+y_x^2)}{y_x} \omega_{xx} - \frac{2y_x y_{xx} \omega_{xx}}{y_x} \\
+ \frac{(1+y_x^2)}{y_x^2} y_{x\psi} \omega_{x\psi}
\]  
(D1)

From the second equation of (2.3.7a), differentiating with respect to \( x \),

\[
R_c h_{\psi x} = -R_c \omega_x + y_x \omega_{xx} + y_{\psi x} \omega_x - y_x \omega_{\psi x} - y_{xx} \omega_{\psi}
\]  
(D2)

Equating (D1) and (D2) and reordering, we get

\[
y_x^2 \omega_{xx} - 2y_x y_{\psi x} \omega_x + (1+y_x^2) \omega_{x\psi} - R_c y_x \omega_x \\
+ 2y_x y_{xx} \omega_{x\psi} - \frac{(1+y_x^2)}{y_x} y_{x\psi} \omega_x - y_{xx} y_{x\psi} \omega_{x\psi} = 0
\]  
(D3)

Now, from (2.3.7b), writing the equation out in full

\[
y_x^2 \omega_{xx} - 2y_x y_x \omega_{xx} + (1+y_x^2) \omega_{xx} - R_c y_x \omega_x - y_x^2 \omega_{x\psi} = 0
\]  
(D4)

(D3) and (D4) will be the same equation if we can show that

\[
- y_x^2 \omega_{x\psi} = 2y_x y_{xx} \omega_{x\psi} - \frac{(1+y_x^2)}{y_x} y_{x\psi} \omega_x - y_{xx} y_{x\psi} \omega_{x\psi}
\]
\[
\text{RHS} = \frac{2y_y y_x y_{\psi} \omega_{\psi} - (1+y_x^2)y_{\psi x} \omega_x - y_{xx} y_{\psi}^2 \omega_x}{y_{\psi}}
\]

\[
- \frac{[y_{\psi}^2 y_{xx} - 2y_x y_{\psi} y_{xx} + (1+y_{\psi}^2)y_{xx}^2] \omega_x}{y_{\psi}}
\]

- \frac{-L\{y\} \omega_{\psi}}{y_{\psi}} \quad \text{from definition of operator L}

- \frac{-y_x^3 \omega_{\psi}}{y_{\psi}} \quad \text{using (2.3.8c)}

- \frac{-y_{\psi}^3 \omega_{\psi}}{}

- \text{LHS}
APPENDIX E

Derivation of the equation for energy $h = h(x, \psi)$, and hence pressure $p = p(x, \psi)$, from equations (2.3.7a)

From definition of the operator $L$,

$L\{h\} = y_x^2 h_{xx} - 2y_x y_{x\psi} h_{x\psi} + (1+y_x^2) h_{\psi\psi}$

\[- \frac{y_x^2}{R_e} \left[ y_x \omega_{xx} + y_{xx} \omega_x - \frac{y_x [ (1+y_x^2) \omega_{x\psi} + 2y_{x\psi} \omega_x ] - (1+y_x^2) \omega_x \omega_{x\psi}}{y_x^2} \right] \]

\[- \frac{2y_x y_{\psi}}{R_e} \left[ y_x \omega_{x\psi} + y_{x\psi} \omega_x - \frac{y_x [ (1+y_x^2) \omega_{x\psi} + 2y_{x\psi} \omega_x ] - (1+y_x^2) \omega_x \omega_{x\psi}}{y_x^2} \right] \]

\[+ (1+y_x^2) \left[ -\omega_x + \frac{1}{R_e} \{ y_x \omega_{x\psi} + y_{x\psi} \omega_x - y_x \omega_{x\psi} - y_{x\psi} \omega_x \} \right] \]

\[- \omega_x \left[ \frac{y_{xx} y_{\psi}^2}{R_e} - \frac{2y_x y_{x\psi} y_{\psi}}{R_e} \frac{(1+y_x^2)}{y_x^2} \right] \]

\[+ \omega_x \left[ -\frac{y_x^2}{R_e} \frac{1}{y_{\psi}} 2y_x y_{x\psi} + \frac{y_x^2 (1+y_x^2)}{R_e} y_{x\psi} - \frac{2y_x y_{\psi}}{R_e} \frac{1}{y_{\psi}} - \frac{1}{R_e} y_{x\psi} \right] \]

\[- \frac{2y_x y_{\psi} (1+y_x^2)}{R_e} \frac{1}{y_{\psi}} - \frac{(1+y_x^2)}{R_e} y_{x\psi} \]

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\[ + y_x \left[ \frac{y_x^2}{R_e} \omega_x \omega_x - \frac{2y_x y_y}{R_e} \omega_y \omega_y - \frac{(1+y_x^2)}{R_e} \omega_{yy} \right] \]

\[ + y_y \left[ \frac{-y_y^2}{R_e} \frac{1}{y_y} (1+y_x^2) \omega_{xy} + \frac{2y_x y_y}{R_e} \frac{1}{y_y} (1+y_x^2) \omega_{yy} + \frac{(1+y_x^2)}{R_e} \omega_{xx} \right] \]

\[ = \frac{\omega_x}{R_e} \left[ y_x^2 \omega_x - 2y_x y_y \omega_{xy} + (1+y_x^2) \omega_{yy} \right] \]

\[ + \omega_y \left[ -(1+y_x^2) \frac{1}{R_e} \left\{ -2y_x y_y \omega_{xx} + 4y_x^2 y_{xy} - \frac{2y_x}{y_y} (1+y_x^2) \omega_{yy} \right\} \right] \]

\[ + \frac{y_x}{R_e} \left[ y_x^2 \omega_{xx} - 2y_x y_y \omega_{xy} - (1+y_x^2) \omega_{yy} \right] \]

\[ + \frac{y_y}{R_e} \frac{(1+y_x^2)}{R_e} \left\{ -\omega_{xx} + \frac{2y_x y_y}{y_y} \omega_{xy} + \omega_{yy} \right\} \]

\[ = \frac{\omega_x y_x^2 \omega_{xx}}{R_e} + \omega_y \left[ -(1+y_x^2) \frac{2y_x}{R_e y_y} \left\{ -y_x^2 \omega_{xx} + 2y_x y_y \omega_{xy} - (1+y_x^2) \omega_{yy} \right\} \right] \]

\[ + \frac{y_x}{R_e} \left[ y_x^2 \omega_{xx} - 2y_x y_y \omega_{xy} + (1+y_x^2) \omega_{yy} \right] \]
\[- \frac{\omega_y y_y \omega}{R_e} + \omega_\phi \left[ -(1 + y_2) - \frac{2y_y y_y \omega}{R_e} \right] + \frac{y_1}{R_e} \left( R_y \omega_x + y_y^2 \omega_\phi \right) \]

\[- \omega_x \left[ \frac{y_2 \omega}{R_e} + y_x y_\phi \right] + \omega_\phi \left[ -(1 + y_2) - \frac{2y_y y_y \omega}{R_e} + \frac{y_1 y_y \omega}{R_e} \right] \]

Hence,

\[ L\{h\} = y_\phi \left[ y_x + \frac{y_2 \omega}{R_e} \right] \omega_x - \left[ 1 + y_2 + \frac{y_2 y_y}{R_e} \omega_\phi \right] \]

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APPENDIX F

Derivation of boundary conditions for vorticity equations (2.4.5.7b) and (2.4.5.7d)

1. \( \omega_{i1} = aq_{1i}^2 + bq_{1i}^2 = -\frac{1}{2}(q_{2i}^2)_{11} \)

Expand in a Taylor Series about \( ii \)

<table>
<thead>
<tr>
<th>( aq_{1i}^2 )</th>
<th>( (q^2)_{11} )</th>
<th>( (q_{1i}^2)_{11} )</th>
<th>( (q_{2i}^2)_{11} )</th>
<th>( (q_{3i}^2)_{11} )</th>
<th>( (q_{4i}^2)_{11} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( a[1(1)] )</td>
<td>1(2A(\psi) )</td>
<td>( \frac{1}{2}(2A(\psi)^2 )</td>
<td>1/6(2A(\psi)^3 )</td>
<td>1/24(2A(\psi)^4 )</td>
<td></td>
</tr>
<tr>
<td>( a )</td>
<td>2A(\psi )</td>
<td>2A(\psi^2 )</td>
<td>(4/3)aA(\psi^3 )</td>
<td>2/3aA(\psi^4 )</td>
<td></td>
</tr>
<tr>
<td>( bq_{1i}^2 )</td>
<td>( b[1(1)] )</td>
<td>1(A(\psi )</td>
<td>( \frac{1}{2}(A(\psi)^2 )</td>
<td>1/6(A(\psi)^3 )</td>
<td>1/24(A(\psi)^4 )</td>
</tr>
<tr>
<td>( b )</td>
<td>bA(\psi )</td>
<td>bA(\psi^2 )</td>
<td>1/6bA(\psi^3 )</td>
<td>1/24bA(\psi^4 )</td>
<td></td>
</tr>
</tbody>
</table>

\( (q_{4i}^2)_{11} \)

\( \Sigma \)

\( a + b \) | \( (2a + b)A\(\psi + \frac{1}{2} \) | \( (2a + b^2)A\(\psi^2 \) | \( (4a/3 + b^2)A\(\psi^3 \) | \( (2a/3 + b/24)A\(\psi^4 \) |
Setting the first two coefficient sums equal to zero gives

\[ a + b = 0 \]

\[ (2a + b)\Delta \psi + \frac{1}{2} = 0 \]

Solving these equations,

\[ b = \frac{1}{2} \Delta \psi \quad a = -\frac{1}{2} \Delta \psi \]

and

\[ \omega_n = -\left(\frac{1}{2} \Delta \psi \right) (q_3^2 - q_2^2) \]

Then, the truncation error \( E_T \) is given by

\[ E_T + (2a + \frac{1}{2} b) \Delta \psi^2 = 0 \]

that is,

\[ E_T = \frac{3}{4} \Delta \psi \]

\[ = 0(\Delta \psi) \]
2. \[ \omega_{11} = aq_{11}^2 + bq_{12}^2 + cq_{11}^2 = -i_4 (q_{xy}^2)_{11} \]

Expand in a Taylor Series about \( il \)

<table>
<thead>
<tr>
<th>( (q^2)_{11} )</th>
<th>( (q_{xy}^2)_{11} )</th>
<th>( (q_{xx}^2)_{11} )</th>
<th>( (q_{yy}^2)_{11} )</th>
<th>( (q_{xy}^2)_{11} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( a[1(1) )</td>
<td>( 1(2\Delta\Psi) )</td>
<td>( \frac{1}{2}(2\Delta\Psi)^2 )</td>
<td>( \frac{1}{6}(2\Delta\Psi)^3 )</td>
<td>( \frac{1}{24}(2\Delta\Psi)^4 )</td>
</tr>
<tr>
<td>( a )</td>
<td>( 2a\Delta\Psi )</td>
<td>( 2a\Delta\Psi^2 )</td>
<td>( \frac{4}{3}a\Delta\Psi^3 )</td>
<td>( \frac{2}{3}a\Delta\Psi^4 )</td>
</tr>
<tr>
<td>( b[1(1) )</td>
<td>( 1(\Delta\Psi) )</td>
<td>( \frac{1}{2}(\Delta\Psi)^2 )</td>
<td>( \frac{1}{6}(\Delta\Psi)^3 )</td>
<td>( \frac{1}{24}(\Delta\Psi)^4 )</td>
</tr>
<tr>
<td>( b )</td>
<td>( b\Delta\Psi )</td>
<td>( \frac{1}{2}b\Delta\Psi^2 )</td>
<td>( \frac{1}{6}b\Delta\Psi^3 )</td>
<td>( \frac{1}{24}b\Delta\Psi^4 )</td>
</tr>
<tr>
<td>( c )</td>
<td>( c )</td>
<td>( c )</td>
<td>( c )</td>
<td>( c )</td>
</tr>
</tbody>
</table>

\[ E = a + b + c \quad (2a+b)\Delta\Psi + \frac{1}{4} \quad (2a+b)\Delta\Psi^2 \quad (4a/3+b/6)\Delta\Psi^3 \quad (2a/3+b/24)\Delta\Psi^4 \]
Set the first three coefficient sums equal to zero:

\[ a + b + c = 0 \]

\[ (2a + b)\Delta\psi + \frac{1}{2} = 0 \]

\[ 2a + \frac{1}{2}b = 0 \]

Hence,

\[ b = -\frac{1}{\Delta\psi} \]

\[ a = \frac{1}{4\Delta\psi} \]

\[ c = \frac{3}{4\Delta\psi} \]

and

\[ \omega_{ii} \approx \left(\frac{1}{4\Delta\psi}\right)\left(q_{ii}^2/4 - q_{ii}^2 + 3q_{ii}^2/4\right) \]

\[ = \left(\frac{1}{4\Delta\psi}\right)\left(q_{ii}^2 - 4q_{ii}^2\right) \text{ when } q_{ii}^2 = 0 \]

As before,

\[ E_T + (4a/3 + b/6)\Delta\psi^3 = 0 \text{ where } E_T \text{ is the truncation error.} \]

Therefore, \[ E_T = -\left[\left(\frac{4}{3}\right)\left(\frac{1}{4\Delta\psi}\right) - \left(\frac{1}{6}\right)\left(\frac{1}{\Delta\psi}\right)\right]\Delta\psi^3 \]

\[ = -\frac{1}{6}\Delta\psi^2 \]

\[ = 0(\Delta\psi^2) \]
APPENDIX G

Derivation of equations in stretched coordinates

From equation (2.4.6.1), for the transformation \( \psi = fn(\eta) \) only, \( \Delta \psi \) transforms to \( \Delta \eta \) and we have

\[
\begin{align*}
\partial / \partial \psi &= (d\psi / d\eta)^{-1}(\partial / \partial \eta) \\
\partial^2 / \partial \psi^2 &= (d\psi / d\eta)^{-2}(\partial^2 / \partial \eta^2) - (d\psi / d\eta)^{-3}(d^2 \psi / d\eta^2)(\partial / \partial \eta) \\
\partial^2 / \partial x \partial \psi &= (d\psi / d\eta)^{-1}(\partial^2 / \partial x \partial \eta)
\end{align*}
\]

\[\text{(G1)}\]

From equations (2.4.1.1) and (2.4.1.2) using (G1) we get

\[
\begin{align*}
(A_{ij}^{(k)} \delta_{xx} + B_{ij}^{(k)} (d\psi / d\eta)^{-1} \delta_{x\eta})
+ C_{ij}^{(k)} [(d\psi / d\eta)^{-2} \delta_{\eta\eta} - (d\psi / d\eta)^{-3} (d^2 \psi / d\eta^2) \delta_{\eta}] \\
+ \alpha R e D_{ij}^{(k)} \delta_x + E_{ij}^{(k)} (d\psi / d\eta)^{-1} \delta_{\eta} \phi_{ij}^{(n+1)} = 0
\end{align*}
\]

Therefore,

\[
\begin{align*}
(A_{ij}^{(k)} \delta_{xx} + B_{ij}^{(k)} (d\psi / d\eta)^{-1} \delta_{x\eta} + C_{ij}^{(k)} (d\psi / d\eta)^{-2} \delta_{\eta\eta} \\
+ \alpha R e D_{ij}^{(k)} \delta_x \\
+ [E_{ij}^{(k)} (d\psi / d\eta)^{-1} - C_{ij}^{(k)} (d\psi / d\eta)^{-3} (d^2 \psi / d\eta^2) \delta_{\eta}] \phi_{ij}^{(n+1)} = 0
\end{align*}
\]

where \( \phi = \begin{cases} y(x, \eta) & \text{if } \alpha = 0 \\ \omega(x, \eta) & \text{if } \alpha = 1 \end{cases} \) are the unknowns,

\[
\begin{align*}
A_{ij}^{(k)} &= (d\psi / d\eta)^{-2} (\delta_{\eta} y)_{ij}^2 \\
B_{ij}^{(k)} &= -2 (\delta_{xy})_{ij} (d\psi / d\eta)^{-1} (\delta_{\eta} y)_{ij} \\
C_{ij}^{(k)} &= 1 + (\delta_{xy})_{ij}^2 \\
D_{ij}^{(k)} &= -(d\psi / d\eta)^{-1} (\delta_{\eta} y)_{ij} \\
E_{ij}^{(k)} &= -(d\psi / d\eta)^{-2} (\delta_{\eta} y)_{ij}^2 \phi_{ij}
\end{align*}
\]

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Simplifying we obtain

\[
(\tilde{A}_{ij}^{(k)} \delta_{xx} + \tilde{B}_{ij}^{(k)} \delta_{x\eta} + \tilde{C}_{ij}^{(k)} \delta_{\eta\eta} + \alpha R_n \tilde{D}_{ij}^{(k)} \delta_x + \tilde{E}_{ij}^{(k)} \delta_\eta) \phi_{ij}^{(n+1)} = 0
\]

where

\[
\tilde{A}_{ij}^{(k)} = (d\psi/d\eta)^{-\frac{1}{2}} (\delta_\eta_1)^{i_1}_1
\]

\[
\tilde{B}_{ij}^{(k)} = -2 (d\psi/d\eta)^{-\frac{1}{2}} (\delta_x^1_1)^{i_1}_1 (\delta_\eta_1)^{i_1}_1
\]

\[
\tilde{C}_{ij}^{(k)} = (d\psi/d\eta)^{-\frac{1}{2}} [1 + (\delta_x^1_1)^{i_1}_1]
\]

\[
\tilde{D}_{ij}^{(k)} = -(d\psi/d\eta)^{-\frac{1}{2}} (\delta_\eta_1)^{i_1}_1
\]

\[
\tilde{E}_{ij}^{(k)} = -(d\psi/d\eta)^{-\frac{1}{2}} (\delta_x^1_1)^{i_1}_1 \omega_{ij} - [1 + (\delta_x^1_1)^{i_1}_1] (d\psi/d\eta)^{-\frac{1}{2}} (d^2\psi/d\eta^2)_1
\]

Therefore, equations (2.4.2.3a) become

\[
\tilde{a}_{ij} = \tilde{A}_{ij}^{(k)}
\]

\[
\tilde{b}_{ij} = \tilde{B}_{ij}^{(k)} (\Delta x/4 \Delta \eta)
\]

\[
\tilde{c}_{ij} = \tilde{C}_{ij}^{(k)} (\Delta x^2/\Delta \eta^2)
\]

\[
\tilde{d}_{ij} = \alpha R_n \tilde{D}_{ij}^{(k)} (\Delta x/2)
\]

\[
\tilde{e}_{ij} = \tilde{E}_{ij}^{(k)} (\Delta x^2/2 \Delta \eta)
\]
From equation (2.4.6.1) \( \Delta x \) transforms to \( \Delta \xi \) and \( \Delta \psi \) transforms to \( \Delta \eta \) and, in addition to the relations in \((G1)\) we have

\[
\frac{\partial}{\partial x} = (dx/d\xi)^{-1}(\partial/\partial \xi)
\]

\[
\frac{\partial^2}{\partial x^2} = (dx/d\xi)^{-2}(\partial^2/\partial \xi^2) - (dx/d\xi)^{-3}(d^2x/d\xi^2)(\partial/\partial \xi)
\] \(\text{(G2)}\)

Again, from equations (2.4.1.1) and (2.4.1.2) using \((G1)\) and \((G2)\) we get

\[
\left[ A_{ij}^{(k)} \left( \frac{dx}{d\xi} \right)^{-2} \delta_{\xi \xi} - \left( \frac{dx}{d\xi} \right)^{-3} \left( \frac{d^2x}{d\xi^2} \right) \delta_{\xi} \right] + B_{ij}^{(k)} \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\xi \eta} + C_{ij}^{(k)} \left[ \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\eta \eta} - \left( \frac{d\psi}{d\eta} \right)^{-3} \left( \frac{d^2\psi}{d\eta^2} \right) \delta_{\eta} \right] + \alpha R \delta_{ij}^{(k)} \left( \frac{dx}{d\xi} \right)^{-1} \delta_{\xi} + E_{ij}^{(k)} \left( \frac{d\psi}{d\eta} \right)^{-3} \delta_{\eta} \right] \Phi_{ij}^{(n+1)} = 0
\]

Therefore,

\[
\left[ A_{ij}^{(k)} \left( \frac{dx}{d\xi} \right)^{-2} \delta_{\xi \xi} + B_{ij}^{(k)} \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\xi \eta} + C_{ij}^{(k)} \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\eta \eta} + \alpha R \delta_{ij}^{(k)} \right] \delta_{\xi} \right] \Phi_{ij}^{(n+1)} = 0
\]

where \( \Phi = \begin{cases} y(\xi, \eta), & \text{if } \alpha = 0 \\ \omega(\xi, \eta), & \text{if } \alpha = 1 \end{cases} \) are the unknowns,

and

\[
A_{ij}^{(k)} = \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\psi \psi}
\]

\[
B_{ij}^{(k)} = -2 \left( \frac{dx}{d\xi} \right)^{-1} \delta_{\xi \xi} \left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\xi \psi} \]

\[
C_{ij}^{(k)} = 1 + \left( \frac{dx}{d\xi} \right)^{-2} \delta_{\xi \xi}
\]

\[
D_{ij}^{(k)} = -\delta_{\psi \eta}
\]

\[
E_{ij}^{(k)} = -\left( \frac{d\psi}{d\eta} \right)^{-2} \delta_{\eta \eta} \omega_{ij}
\]

Simplifying we obtain

\[
\left( \overline{A}_{ij}^{(k)} \right) \delta_{\xi \xi} + \overline{B}_{ij}^{(k)} \delta_{\xi \eta} + \overline{C}_{ij}^{(k)} \delta_{\eta \eta} + \overline{D}_{ij}^{(k)} \delta_{\xi} + \overline{E}_{ij}^{(k)} \delta_{\eta} \right] \Phi_{ij}^{(n+1)} = 0
\]
where
\[ \overline{A}_{ij}^{(k)} = \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \]
\[ \overline{B}_{ij}^{(k)} = -2 \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_x Y \right)_{ij} \left( \delta_n Y \right)_{ij} \]
\[ \overline{C}_{ij}^{(k)} = \left( \frac{d\psi}{d\eta} \right)^{-1} \left[ 1 + \left( \frac{dx}{d\xi} \right)^{-1} \left( \delta_x Y \right)_{ij} \right] \]
\[ \overline{D}_{ij}^{(k)} = -\alpha R_n \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \left( \frac{dx}{d\xi} \right)^{-1} \]
\[ -\left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{dx}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d^2x}{d\xi^2} \right) \]
\[ = -\left( \frac{dx}{d\xi} \right)^{-1} \left( \alpha R_n \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \right) \]
\[ + \left( \frac{dx}{d\xi} \right)^{-1} \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \frac{dx}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \left( \frac{d^2x}{d\xi^2} \right) \]
\[ \overline{E}_{ij}^{(k)} = -\left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \omega_{ij} \]
\[ -\left[ 1 + \left( \frac{dx}{d\xi} \right)^{-1} \left( \delta_x Y \right)_{ij} \left( \frac{d\psi}{d\eta} \right)^{-1} \left( \frac{d^2\psi}{d\eta^2} \right) \right] \]
\[ = -\left( \frac{d\psi}{d\eta} \right)^{-1} \left( \delta_n Y \right)_{ij} \omega_{ij} + \left( \frac{d\psi}{d\eta} \right)^{-1} \left[ 1 + \left( \frac{dx}{d\xi} \right)^{-1} \left( \delta_x Y \right)_{ij} \right] \]

Therefore, equations (2.4.2.3a) become
\[ \overline{a}_{ij} = \overline{A}_{ij}^{(k)} \]
\[ \overline{b}_{ij} = \overline{B}_{ij}^{(k)} \left( \Delta \xi / 4 \Delta \eta \right) \]
\[ \overline{c}_{ij} = \overline{C}_{ij}^{(k)} \left( \Delta \xi ^2 / \Delta \eta \right) \]
\[ \overline{d}_{ij} = \overline{D}_{ij}^{(k)} \left( \Delta \xi / 2 \right) \]
\[ \overline{e}_{ij} = \overline{E}_{ij}^{(k)} \left( \Delta \xi ^2 / 2 \Delta \eta \right) \]
APPENDIX H

Channel shape

The channel shape for this study was chosen by Roache [III.1]. It was desirable to have separated flow with a nearly constant channel height at outflow, and a shape defined by a single-valued smooth analytic function. The shape of the lower wall was defined by a shifted hyperbolic tangent function, given by

\[ y - y_f(x) = \frac{1}{2} \left[ 1 - \tanh \left( sc \frac{x - x_F}{L_c} \right) \right] + \delta \]  \hspace{1cm} (H1)

where \( L_c \) = length of the channel;

\( x_F \) = location of the inflection point in the \( \tanh \) profile;

\( sc \) = scale factor;

\( \delta \) = small adjustment set to give \( y_f(0) = 0 \).

For the parametric cases considered, \( L_c \) is scaled with Reynolds number \( R_c = R_{\infty} \) as

\[ L_c = \frac{R_{\infty}}{C_{t_{\infty}}} \]. Again, this scaling is necessary to keep the separated flow region within the computational mesh, and results in the self-similarity of the solutions at high \( R_c \). The particular channel parameters used were:

\[ sc = 10, \quad \frac{x_F}{L_c} = 0.2 \quad \text{so that} \quad x_F = \frac{R_{\infty}}{15} \quad \text{and} \quad C_{t_{\infty}} = 3 \]  \hspace{1cm} (H2)
Substituting (H2) into (H1) and taking \( y_t(0) = 0 \), we obtain

\[
y_t(x) = \frac{1}{2} \left[ 1 - \tanh \left( \frac{10 \frac{x-x_F}{R_\infty}}{3} \right) \right] + \delta
\]

\[
= \frac{1}{2} \left[ \tanh \left( 2 - \frac{30x}{R_\infty} \right) + 1 \right] + \delta
\]

\[
= \frac{1}{2} \left[ \tanh \left( 2 - \frac{30x}{R_\infty} \right) - \tanh 2 \right]
\]
APPENDIX I

Justification of the boundary conditions at the channel outlet

The slope of the lower wall at is given by

\[
y' = \frac{1}{2} \text{sech}^2 \left( 2 - \frac{30x}{R_e} \right) \left[ -\frac{30}{R_e} \right]
\]

Hence, at \( x = x_{\text{out}} = \frac{R_e}{3} \),

\[
y' \left( \frac{R_e}{3} \right) = \frac{-15}{R_e} \text{sech}^2 (2 - 10)
\]

\[
= 7 \times 10^{-7} \text{ for } R_e = 10
\]

\[
= 7 \times 10^{-4} \text{ for } R_e = 100
\]

The channel shape has a height that is constant to five figures over the last 20% of the channel. Hence, using the condition commonly used in straight-channel calculations, i.e., parabolic velocity profile, would be acceptable. In any case, the lower wall is virtually parallel to the x-axis at the outlet, and we can reasonably assume that \( v = 0 \) at the outlet. To further support the above argument, consider the following:

\[
u = \psi_y
\]

\[
v = -\psi_x
\]

Hence, \( u_x = 0 \) implies \( \psi_{yx} = 0 \), i.e., \( (\psi_y)_y = 0 \). Similarly, \( v_x = 0 \) implies \( -\psi_{xx} = 0 \), i.e., \( (-\psi_x)_x = 0 \).

Therefore, we can conclude that \( \psi_x = -v = \text{constant} \). Since \( v = 0 \) at the walls, this
constant is zero, and hence \( v = 0 \) at \( x = x_{\text{out}} \) for all \( y \). Also, \( v = uy_x = \frac{y_x}{y_y} = 0 \)

implies \( y_x = 0 \) for \( x = x_{\text{out}} \) and \( y_f \left( \frac{R_{\infty}}{3} \right) \leq y \leq 1 \).

Finally, \( \omega = v_x - u_y \), so

\[
\omega \big|_{x=x_{\text{out}}} = -u_y = -\psi_{yy}
\]

Hence,

\[
\omega_x \big|_{x=x_{\text{out}}} = -\psi_{yy} = (-\psi_{yx})_x = 0
\]

Therefore, \( \omega_x = 0 \) for \( x = x_{\text{out}} \) and \( y_f \left( \frac{R_{\infty}}{3} \right) \leq y \leq 1 \).
APPENDIX J

Derivation of equation (3.2.3.1)

We want to solve \( \psi = \frac{1}{2} (3y^2 - y^3) \) explicitly for \( y = y(\psi) \). Rewrite this equation as

\[
y^3 - 3y^2 + 2\psi = 0 \tag{J1}
\]

Let \( p = -3, \ y = x - \frac{p}{3} = x + 1 \). Then

\[
(x + 1)^3 - 3(x + 1)^2 = 2\psi = 0
\]

or,

\[
x^3 - 3x - 2 + 2\psi = 0 \tag{J2}
\]

Let \( a = -3, \ b = -2 + 2\psi \), then \( ab = -3(-2 + 2\psi) \neq 0 \) provided \( \psi \neq 1 \).

The solution of (J2) is given by

\[
x = m \cos \theta
\]

where

\[
m = 2 \sqrt{\frac{-a}{3}} - 2
\]

and

\[
\cos 3\theta = \frac{3b}{9m} - 1 - \psi
\]

Hence, the solution of (J1) is

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\[ y = m \cos \theta + 1 \]
\[ = 2 \cos \theta + 1 \]  

where

\[
\theta = \begin{cases} 
\theta_1, \\
\theta_1 + \frac{2\pi}{3}, \\
\theta_1 + \frac{4\pi}{3}
\end{cases}
\]  

and \( \theta_1 \) is defined as

\[
\theta_1 = \frac{1}{3} \cos^{-1}(1 - \psi)
\]

Several possibilities occur for the value of \( \theta \). To determine the appropriate choice, consider the following cases:

1. When \( \psi = 0 \), we require that \( y = 0 \). From (I1), \( \psi = 0 \) yields

\[ y^2(y-3) = 0 \]

This implies that \( y = 0 \) or \( y = 3 \), so we can select \( y = 0 \). Using \( \psi = 0 \) in (I5) gives

\[ \theta_1 = \frac{1}{3} \cos^{-1}(1) = 0 \]
Then, (J4) gives

\[
\theta = \begin{cases} 
0 \\
\frac{2\pi}{3} \\
\frac{4\pi}{3} \\
\frac{\pi}{3}
\end{cases}
\]

We find that \(y = 0\) for \(\theta = \frac{2\pi}{3}\) and \(\frac{4\pi}{3}\).

2. Equation (J1) is also satisfied at the channel centreline, \(y = 1\), at which \(\psi = 1\).

Hence, using (J5)

\[
\theta_1 = \frac{1}{3} \cos^{-1}(0) = \frac{\pi}{6}
\]

Using (J4),

\[
\theta = \begin{cases} 
\frac{\pi}{6} \\
\frac{\pi}{6} + \frac{2\pi}{3} \\
\frac{\pi}{6} + \frac{4\pi}{3} \\
\frac{\pi}{3}
\end{cases}
\]

We find that \(y = 1\) for \(\theta = \frac{\pi}{6} + \frac{4\pi}{3}\) only.

Combining these results, we see that the only choice for \(\theta\) is \(\theta = \theta_1 + \frac{4\pi}{3}\).

Hence, we arrive at the solution

\[
y = 2\cos \left( \theta_1 + \frac{4\pi}{3} \right) + 1
\]

where

\[
\theta_1 = \frac{1}{3} \cos^{-1}(1 - \psi)
\]
APPENDIX K

Difference formulas used for $v_x$ in equation (3.2.4.1)

We have $u = \frac{1}{y_\phi}$ and $v = uy_x = \frac{y_x}{y_\phi}$.

The derivative $v_x$ is approximated by forward differencing,

$$(v_x)_i \approx (\delta_x v)_i = \frac{v_{i+1,j} - v_i}{\Delta x} \quad 2\text{-point, } O(\Delta x)$$

or

$$\frac{(\delta_x v)_i}{2} = \frac{-v_{i+2,j} + 4v_{i+1,j} - 3v_i}{2\Delta x} \quad 3\text{-point, } O(\Delta x)$$

We use a forward (downwind) difference for $y_x$, but a central difference for $y_\phi$, to compute values of $v$,

$$v_i = \begin{bmatrix} \frac{\delta_x y}{\delta y} \\ \frac{\delta y}{\delta y} \\ \end{bmatrix} = \frac{y_{i+1,j} - y_{i,j}}{\Delta x} \frac{2\Delta \psi}{y_{i+1,j} - y_{i,j-1}}, \quad 0(\Delta x)0(\Delta \psi)$$

or

$$v_i = \frac{-y_{i,2,j} + 4y_{i+1,j} - 3y_{i,j}}{2\Delta x} \frac{\Delta \psi^2}{y_{i,j+1} - 2y_i + y_{i,j-1}}, \quad 0(\Delta x)0(\Delta \psi^2)$$

The above formulas are valid for $i = 1$ only, $j = 2,3,\ldots,J_1$.

At $j = 1$, the solid wall no-slip condition gives $(v_x)_{1,j} = 0$.

At $j = J_1$, the centreline symmetry condition implies $(v_x)_{j,J_1} = 0$. 

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APPENDIX L

Derivation of the solution for a circular cylinder in hyperbolic-cosine shear flow

Problem A: Flow about a circular cylinder of radius \( r = a \)

Consider steady plane (2-dimensional) motion of an incompressible, inviscid fluid past a circular cylinder without circulation (irrotational).

\[ \nabla^2 \psi = -\omega(\psi) = 0, \text{ since the flow is irrotational.} \]

\( \psi \) is the stream function with velocity components in cartesian coordinates given by \( u = \psi_y, \ v = -\psi_x, \) and \( \omega = \omega(\psi) \) is the vorticity.

**BC's:**

Upstream: Uniform flow far upstream at infinity, implying that the vorticity vanishes.

\[
\begin{align*}
\text{Speed:} & \quad \begin{cases}
\text{at infinity.} \\
u = U_\infty \quad \text{constant} \\
v = 0
\end{cases}
\end{align*}
\]

Stream function: Integrating \( \psi_y = U_\infty \) gives

\[ \psi_\infty = U_\infty y = U_\infty rsin\theta \text{ in polar coordinates } (r, \theta). \] Hence

\[ \psi(r, \theta) \to \psi_\infty = U_\infty rsin\theta \text{ as } r \to \infty. \]

Note: We have taken \( y = 0 \) as the streamline \( \psi = 0 \) so the arbitrary constant of integration will be zero. This will be the case in the problems to follow.

Surface: Since the surface is a streamline, we can take \( \psi = 0 \) on \( r = a, \) i.e., \( \psi(a, \theta) = 0. \)
Circulation: To rule out circulation and obtain a unique solution, it is assumed that no flow will cross the x axis by assuming symmetry about the line \( \theta = 0 \) and \( \theta = \pi \) \((|r| > a)\), i.e., no additional circulation is induced by the body.

Solution: \[
\psi(r, \theta) = U \left[ r - \frac{a^2}{r} \right] \sin \theta
\]
i.e., uniform stream plus a dipole at the centre of the circle.
Problem B: Circular cylinder in slight shear flow (same as problem A except as noted below)

DE: \( \nabla^2 \psi = -\omega(\psi) \neq 0, \) since the flow is rotational. \( \text{(L1)} \)

BC's: Upstream: Consider a slight or small linear perturbation to the uniform flow boundary condition far upstream at infinity. Let the oncoming stream be a parallel flow with small constant vorticity.

Speed:
\[
\begin{align*}
  u &= U_\infty \left( 1 + \epsilon \frac{y}{a} \right) \\
  v &= 0 \\
\end{align*}
\]

at infinity.

Stream function: Integrating \( \psi_{\infty y} = U_\infty \left( 1 + \epsilon \frac{y}{a} \right) \) gives

\[
\psi_\infty = U_\infty \left( y + \frac{1}{2} \epsilon \frac{y^2}{a} \right)
\]

That is, \( \psi(r, \theta; \epsilon) \to \psi_\infty - U_\infty \left[ r \sin \theta + \frac{1}{4} \epsilon \frac{r^2}{a} (1 - \cos 2\theta) \right] \)

as \( r \to \infty. \)
Vorticity: \[ \omega_{\infty} = -\nabla^2 \psi_{\infty} = -\frac{\epsilon U_{\infty}}{a} = \text{constant}. \]

Hence, \[ \omega(\psi) = \omega_{\infty} = -\frac{\epsilon U_{\infty}}{a}. \]

DE: Equation (L1) becomes

\[ \nabla^2 \psi = \frac{\epsilon U_{\infty}}{a} \]

Solution: If the dimensionless "vorticity number" \( \epsilon \) is small (slight vorticity), it seems likely that the flow will depart only slightly from the solution in problem A for irrotational motion.

Hence, by perturbing the solution to problem A, we get

\[ \psi(r, \theta; \epsilon) = U_{\infty} \left[ r - \frac{a^2}{r} \right] \sin \theta + \frac{1}{4} \epsilon U_{\infty} \left[ \frac{r^2}{a} (1 - \cos 2\theta) + \frac{a^3}{r^2} \cos 2\theta - a \right] \]

The basic solution consists of a uniform stream (a dipole at infinity) plus its image in the circle (a dipole at the origin) as before. The first order perturbation solution consists of the rotational part of the stream, its image in the circle and a constant to adjust the stream function to zero on the surface.

Note: Since the vorticity is constant everywhere, this problem is not difficult to solve.
Problem C: Circular cylinder in parabolic and hyperbolic-cosine shear (same as problem A except as noted below)

BC's: Upstream: Consider a circular cylinder of radius a symmetrically placed in a parallel stream of incompressible, inviscid fluid having the following velocity profiles far upstream.

a) Parabolic velocity profile

Speed: \[ u = U_\infty \left( 1 + \frac{1}{2} \epsilon \frac{y^2}{a^2} \right) \] at infinity.
\[ v = 0 \]

Stream Function: Integrating \( \psi_\infty - U_\infty \left( 1 + \frac{1}{2} \epsilon \frac{y^2}{a^2} \right) \) and choosing \( \psi = 0 \) along \( y = 0 \) gives
\[ \psi_\infty = U_\infty \left( y + \frac{1}{6} \epsilon \frac{y^3}{a^2} \right) \]

That is, \( \psi(r, \theta; \epsilon) \rightarrow \psi_\infty = U_\infty \left( r \sin \theta + \frac{1}{6} \epsilon \frac{r^3 \sin^3 \theta}{a^2} \right) \)
as \( r \rightarrow \infty \).

Vorticity: \[ \omega_\infty = -\nabla^2 \psi_\infty = -\frac{\epsilon}{a^2} U_\infty y \neq \text{constant.} \]

In inviscid motion we have the physical fact that vorticity is constant along streamlines in the absence of viscosity.

Hence, writing \( \omega(\psi) = \omega = \omega_\infty \)
and \( \psi(r, \theta; \epsilon) = \psi = \psi_\infty \)

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we have \( y = \frac{a^2}{\varepsilon} \frac{\omega}{U_\infty} \) and, therefore,

\[
\psi - U_\infty \left\{ - \frac{a^2}{\varepsilon} \frac{\omega}{U_\infty} - \frac{1}{6} \frac{a^4}{\varepsilon^2} \frac{\omega^3}{U_\infty^2} \right\} \\
- \frac{a^2}{\varepsilon} \left\{ - \omega - 1 \frac{a^2}{6 \varepsilon U_\infty^2} \right\}
\]

Reverting this series gives

\[
-\omega = \omega(\psi) = \frac{\varepsilon}{a^2} \psi - \frac{1}{6} \frac{\varepsilon^2}{U_\infty^2 a^4} \psi^3 + O(\varepsilon^3)
\]

**D.E.:** Equation (L1) now becomes

\[
\nabla^2 \psi = \frac{\varepsilon}{a^2} \psi + \frac{1}{6} \frac{\varepsilon^2}{a^4 U_\infty^2} \psi^3 + O(\varepsilon^3)
\]

**Solution:**

\[
\psi(r, \theta; \varepsilon) = U_\infty \left\{ r \frac{a^2}{r} \right\} \sin \theta + \varepsilon U_\infty \left\{ \frac{1}{6} \frac{r^3}{a^2} \sin^3 \theta - \frac{1}{2} \frac{r \ell n \frac{r}{a}}{\varepsilon} \sin \theta + \chi \right\} + O(\varepsilon^3)
\]

where \( \chi \) is a solution of the homogeneous equation (complementary solution) that restores the boundary conditions. However, the term in the particular integral or solution that contains a logarithm gives velocity perturbations that are logarithmically infinite at infinity and no harmonic function \( \chi \) will cancel them, i.e., no solution exists with disturbances dying out to satisfy the upstream condition.

Hence, there is no analytic solution here to verify numerically.

Therefore, we consider a profile which is almost the same near the cylinder.
b) Hyperbolic-cosine profile

Speed: \[ u = U_\infty \cosh \left( \frac{1}{\epsilon^2} \frac{y}{a} \right) \] at infinity.
\[ v = 0 \]

Stream function: Integrating \( \psi = U_\infty \cosh \left( \frac{1}{\epsilon^2} \frac{y}{a} \right) \) gives

\[ \psi_\infty = \frac{U_\infty a}{\epsilon^2} \sinh \left( \frac{1}{\epsilon^2} \frac{y}{a} \right) \]
\[ \psi_\infty = U_\infty \left( y + \frac{1}{6} \frac{\epsilon y^3}{a^2} + \ldots \right) \]

Thus,

\[ \psi(r, \theta; \epsilon) \to \psi_\infty = U_\infty \left( r \sin \theta + \frac{1}{6} \frac{\epsilon r^3 \sin^3 \theta}{a^2} + \ldots \right) \]
as \( r \to \infty \).

Vorticity: \[ \omega = -\nabla^2 \psi_\infty = -\frac{\epsilon}{a^2} \psi_\infty \neq \text{constant} \]

Hence, writing \( \omega(\psi) = \omega \equiv \omega_\infty \) and

\[ \psi(r, \theta; \epsilon) = \psi = \psi_\infty \], we have

\[ -\omega = -\omega(\psi) = \frac{\epsilon}{a^2} \psi \]

DE: Equation (L1) now becomes exactly

\[ \nabla^2 \psi = \frac{\epsilon}{a^2} \psi \]
Solution: \( \psi(r, \theta; \epsilon) = U_\infty \left( r - \frac{a^2}{r} \right) \sin \theta + \epsilon U_\infty \left( \frac{1}{6} \frac{r_1^4}{a^6} \sin^3 \theta - \frac{1}{2} \frac{r_1 r}{a} \sin \theta + \chi \right) + O(\epsilon^2) \) (L2)

where a solution \( \chi \) can now be found that disturbs the distant flow upstream as little as possible.

We find,

\[
\chi(r, \theta) = -\frac{1}{8} \frac{a^2}{r} \sin \theta + \frac{1}{24} \frac{a^4}{r^3} \sin 3 \theta + c \left( r - \frac{a^2}{r} \right) \sin \theta
\]

where \( c = \frac{1}{4} \left( \frac{\ln \frac{4}{\epsilon} - 2\gamma + 1}{\epsilon} \right) \); \( \gamma = 0.577 \)

Speed:

\[
q^2 = q_r^2 + q_\theta^2 = \frac{1}{r^2} \left( \frac{\partial \psi}{\partial \theta} \right)^2 + \left( \frac{\partial \psi}{\partial r} \right)^2
\]

From equation (L2)

\[
\frac{\partial \psi}{\partial \theta} = U_\infty \left( r - \frac{a^2}{r} \right) \cos \theta + \epsilon U_\infty \left( \frac{1}{2} \frac{r_1^3}{a^2} \sin^2 \theta \cos \theta - \frac{1}{2} \frac{r_1 r}{a} \sin \theta + \frac{\partial \chi}{\partial \theta} \right) + O(\epsilon^2)
\]

\[
\frac{\partial \chi}{\partial \theta} = -\frac{1}{8} \frac{a^2}{r} \cos \theta + \frac{1}{8} \frac{a^4}{r^3} \cos 3 \theta + c \left( r - \frac{a^2}{r} \right) \cos \theta
\]

\[
\frac{\partial \psi}{\partial r} = U_\infty \left( 1 + \frac{a^2}{r^2} \right) \sin \theta + \epsilon U_\infty \left( \frac{1}{2} \frac{r_1^2}{a^2} \sin^2 \theta - \frac{1}{2} \frac{r_1 r}{a} \sin \theta - \frac{1}{2} \sin \theta + \frac{\partial \chi}{\partial r} \right) + O(\epsilon^2)
\]

\[
\frac{\partial \chi}{\partial r} = -\frac{1}{8} \frac{a^2}{r^2} \sin \theta - \frac{1}{8} \frac{a^4}{r^4} \sin 3 \theta + c \left( 1 + \frac{a^2}{r^2} \right) \sin \theta
\]
In order to obtain expressions in nondimensional variables, we can take $\text{U}_\infty = a = 1$.

Then, $u = \cosh(\varepsilon^2 y)$

$$\psi_\infty = \frac{1}{\varepsilon^2} \sinh\left(\frac{1}{\varepsilon^2} y\right) \quad \text{or} \quad y = \frac{1}{\varepsilon^2} \sinh^{-1}\left(\frac{1}{\varepsilon^2} \psi_\infty\right) \quad (L3)$$

$$\omega = -\varepsilon \psi = -\varepsilon^2 \sinh\left(\frac{1}{\varepsilon^2} y\right) \quad (L4)$$

The preceding solution is obtained with considerable difficulty using the method of matched asymptotic expansions. A perturbation solution is required, one valid far away from the body, termed the outer solution and another valid near the body, termed the inner solution. Both of these solutions contain arbitrary constants which have to be matched using some form of matching principle to obtain a unique solution. A uniformly valid composite approximation can be obtained by finding the common part of the inner and outer solutions and subtracting it from the inner plus outer solution using additive composition.
Derivation of the initial guess for equation (5.2.4.4)

We have \( 2 = x_{np} < x < x_A = 2 + L_x \)

For an ellipse we take \( a = L_x \), \( b = h = 0.4 \) to get the equation

\[
\frac{(x-2)^2}{L_x^2} + \frac{y^2}{h^2} = 1
\]

Solving for \( y \) gives

\[
y = h \left[ 1 - \frac{(x-2)^2}{L_x^2} \right]^{\frac{1}{2}} = y_{ii}^{0}\]

in equation (5.2.4.4). This gives the initial guess for \( y \), for some initial value of the reattachment length \( L_x \).
APPENDIX N

Computer program

PROGRAM SLOR

IMPLICIT DOUBLE PRECISION (A-H, O-Z)

EXTERNAL DELXY, DELPSY

PARAMETER (IMAX=61, JMAX=21, IMM = IMAX - 1, JMM1 =
# JMAX - 1, IMM2 = IMAX - 2, JMM2 = JMAX - 2)

DIMENSION A(JMM1), B(JMM1), C(JMM1), D(JMM1), R(JMM1),
# WOLD(IMAX,JMAX), WNEW(IMAX,JMAX), BENCH(JMAX,2),
# YOLD(IMAX,JMAX), YNEW(IMAX,JMAX)

,C, F(JMAX,2), G(JMAX,2)

COMMON /BND/DELX, DELPSY, XMN, PSIMIN, REY

COMMON /WVALUE/WOLD, WNEW/YVALUE/YOLD,

YNEW/WCH/WCHECK/WA/HSTEP

LOGICAL CHECK, WCHECK

C

DEFINITION OF STATEMENT FUNCTIONS

A1(I,J) = DELPSY(I,J) ** 2
A2(I,J) = -2.0 * DELXY(I,J) * DELPSY(I,J)
A3(I,J) = 1.0 + DELXY(I,J) ** 2
A4(I,J) = -DELPSY(I,J)
A5(I,J) = -A1(I,J) * WOLD(I,J)
Q(I,J) = A3(I,J) / A1(I,J)
VK(I) = (YOLD(I+1,1)-YOLD(I-1,1)) / (YOLD(I,2)-YOLD(I,1)) *
# DELPSI * 0.5 / DELX

C

FORMULA 1 OF W

W1(I) = (3.0 * Q(I, 4) - 8.0 * Q(I, 3)
# + 5.0 * Q(I, 2)) / (4.0 * DELPSI)
W2(I) = -(3.0 * Q(I, JMAX-3) - 8.0 * Q(I, JMM2)
# + 5.0 * Q(I, JMM1)) / (4.0 * DELPSI)
W3(I) = -(2.0 * Q(I, 4) - 9.0 * Q(I, 3)
# + 18.0 * Q(I, 2) - 11.0 * Q(I, 1) ) / (12.0 *
DELPSI)

FORMULA 2 OF W

W1(I) = -(2.0 * Q(I, 4) - 9.0 * Q(I, 3)
# + 18.0 * Q(I, 2)) / (12.0 * DELPSI)
W2(I) = (2.0 * Q(I, JMAX-3) - 9.0 * Q(I, JMM2)
# + 18.0 * Q(I, JMM1)) / (12.0 * DELPSI)
W3(I) = 0.5* ((VK(I+1) - VK(I-1)) / DELX -
# (Q(I,2) - Q(I,1)) / DELPSI )

FORMULA 3 OF W

W(I) = -(11.0 * Q(I,5) - 42.0 * Q(I, 4) + 57.0 * Q(I, 3)
# - 26.0 * Q(I, 2)) / (12.0 * DELPSI)

READ THE MINIMUM AND MAXIMUM X AND PSI VALUES. READ THE
VALUE OF EPSILON TO BE USED

WRITE(*,*) ' Input Reynolds Number'
READ(*,*)REY
WRITE(*,*) ' Input Beta for y'
READ(*,*)BETAY

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WRITE(*,*) 'Input Beta for w'
READ(*,*) BETA_W
WRITE(*,*) 'Input delta for w'
READ(*,*) DELTA
WRITE(*,*) 'Input epsilon for error tolerance'
READ(*,*) EPSILY, EPSILW
WRITE(*,*) 'Input maximum iteration Number'
READ(*,*) NMAX
WRITE(*,*) 'Input starting point for delta calculation'
READ(*,*) MMU
WRITE(*,*) 'Input index for reattachment'
READ(*,*) IRETAG, ISTEP
read(*,*) CF
OPEN(UNIT = 1, FILE = 'TEST.DAT', STATUS = 'NEW')

DATA INITIALIZATION

DO 10 I=1,IMAX
   DO 10 J=1,JMAX
      WOLD(I,J) = 0.0
      YOLD(I,J) = 0.0
   CONTINUE
10 CONTINUE

CALCULATE DELTA X AND DELTA PSI

HSTEP = 0.4

ASSUMING AN INDEX FOR REATTACHMENT, WHICH IS CLOSE TO THE OUTLET

XMIN = 0.0
XMAX = 6.0
PSIMIN = 0.0
PSIMAX = 1.0
DEXL = (XMAX - XMIN) / IMM1
DELPSI = (PSIMAX - PSIMIN) / JMM1
RAT = DELX / DELPSI
RAT2 = RAT * RAT
RATD = RAT * DELX * 0.5
RATF = RAT * 0.25
REY = 0.5 * REY * DELX

WRITE(1,111)XMIN, XMAX, PSIMIN, PSIMAX, REY,
          DELTA, BETAY, BETA_W, EPSILY, EPSILW
111 '========================================================================'
# 4X,'Xmin = ',F5.2,' Xmax = ',F5.2/
# 4X,'Pmin = ',F5.2,' Pmax = ',F5.2/
# 4X,'Re = ',F5.2,' delta= ',F5.2/
# 4X,'By = ',F5.2,' Bw = ',F5.2/
# 4X,'Ey = ',E12.5,' Ew = ',E12.5/
# 1X, '========================================================================'

C CALCULATE THE BOUNDARY CONDITIONS WHICH REMAIN UNCHANGED DURING ITERATION
C THEY ARE I<>IMAX FOR Y AND I=1 AND J=JMAX FOR W
C BOUNDARY CONDITIONS FOR Y ON THE LOWER AND UPPER BOUNDARIES
C
DO 20 I=1, IMAX
   YOLD(I,1) = YBND(I,1)
   YOLD(I,IMAX) = YBND(I,IMAX)
   WOLD(I,1) = WBND(I,1)
   WOLD(I,IMAX) = WBND(I,IMAX)
20 CONTINUE
C
C BOUNDARY CONDITIONS OF THE INLET AND OUTLET
C
DO 30 J=1, JMAX
   YOLD(1,J) = YBND(1,J)
   YOLD(IMAX,J) = YBND(IMAX, J)
   WOLD(1,J) = WBND(1,J)
   WOLD(IMAX,J) = WBND(IMAX, J)
30 CONTINUE
C
DO 33 J = 2, JMM1
C
DO 33 I=2, IMM1
   WOLD(I,J) = WOLD(1,J) + (WOLD(IMAX,J) - WOLD(1,J)) *
                FLOAT(I-1) / 60.0
33 CONTINUE
C
C DO ITERATIONS
C
C
DO 32 I=21,29
C32 WRITE(1,*)K,I,YOLD(I,1)
   DO 899 K = 1, NMAX, ISTEP
C
   SET OR RESET THE FLAG IF THE VORTICITY IS COMPUTED
C
   WCHECK = MOD(K,2) .EQ. 0
C
C CALCULATE THE BOUNDARY CONDITIONS FOR W ON LOWER AND UPPER
C BOUNDARIES
C (J=1 OR J=JMAX), WHICH CHANGES DURING THE ITERATIONS
C
   IF(WCHECK) THEN
      IF (K.EQ.2) THEN
         DO 31 J=1,JMAX
            WOLD(I,J) = WBND(I,J)
         31 WOLD(IMAX,J) = WBND(IMAX, J)
      ENDIF
      DO 40 I = 2, IMM1
         WRITE(*,*)YOLD(I,2),YOLD(I,3),YOLD(I,4)
         WRITE(*,*)YOLD(I,IMAX-3),YOLD(I,JMM2),YOLD(I,JMM1)
         IF ( I .GT. 21 .AND. I .LT. IRETMACH ) THEN
            WOLD(I, 1) = (1.0 - DELTA) * WOLD(I,1) + DELTA *
         W3(I)
         ELSE
            WOLD(I, 1) = (1.0 - DELTA) * WOLD(I,1) + DELTA *
         W1(I)
C ENDIF
C DELTA*W2(I)
C WRITE(*,'(I3,F10.4)')I,W1(I),W2(I),W3(I)
C WOLD(I,1) = WBND(I,1)
C WOLD(I,JMAX) = WBND(I,JMAX)
C WRITE(*,*)I,WOLD(I,1),WOLD(I,JMAX)
C CONTINUE
C WRITE(*,'(A)')' OK'
ELSE
  DO 32 J=1,JMAX
    YOLD(IMAX,J) = YBND(IMAX, J)
    WRITE( 1, *)'K=', K
  ENDIF
C DO 41 I = 1, IMAX
C WRITE( *,'(4F10.4)') YOLD(I,1), YOLD(I, JMAX), WOLD(I,1),
C # WOLD(I, JMAX)
C CONTINUE
C WRITE(*,'(A)')'***'
C DO 42 J=1,JMAX
C WRITE(*,'(4F10.4)') YOLD(I,J),YOLD(IMAX,J), WOLD(I,J),
C # WOLD(IMAX,J)
C CONTINUE
C DO 300 I = 2, IMM1
C IF ( K .EQ. 1 .AND. I.GT.21 .AND. I.LT.IRETACH) THEN
C YOLD(I,1) = YOLD(I-1,1)
C ENDIF
C DO 100 J = 2, JMM1
C CALCULATE THE VALUE NEEDED IN THE PRESENT ROW OF THE
C COEFFICIENT MATRIX
C
C1 = A1(I,J)
C2 = A2(I,J) * RATF
C3 = A3(I,J) * RAT2
IF(WCHECK) THEN
  C4 = A4(I,J) * REY2
ELSE
  C4 = 0.0
END IF
C5 = A5(I,J) * RATD
C WRITE(*,'(I5,F10.5)')I, C1,C2,C3,C4,C5
C CALCULATE THE COEFFICIENT MATRIX
C IF(J .GT. 2) THEN
A(J) = C3 - C5
ENDIF
B(J) = -2.0 * (C1 + C3)
IF(J .LT. JMM1) THEN
  C(J) = C3 + C5
ENDIF
FIND THE SOLUTION FOR THE PRESENT I-TH LINE

WRITE(*,'(I5, 4F10.3)') (Z, A(J), B(J), C(J), D(J), J=2,JMM1)
CALL TRID(A,B,C,D,R,2,JMM1)
WRITE(1,*) (R(J), J=1,JMAX)

UPDATE NEW VALUES

IF(WCHECK) THEN
   DO 200 J = 2, JMM1
      WNEW(I,J) = (1.0 - BETAW) * WOLD(I,J) + BETAW * R(J)
   WRITE(1,*) WNEW(I,J)
   CONTINUE
   WRITE(1,*) WNEW(I,2)
ELSE
   DO 201 J = 2, JMM1
      YNEW(I,J) = (1.0 - BETAY) * YOLD(I,J) + BETAY * R(J)
   WRITE(1,*) YNEW(I,J)
   CONTINUE
   WRITE(1,*) YNEW(I,2)
END

MODIFY THE BOUNDARY CONDITIONS FOR Y ON LOWER BOUNDARY (J=1)

IF (I .GT. 21 .AND. I .LE. IRETACH ) THEN
   WRITE(*,*) I, YOLD(I,1), YNEW(I,2), YNEW(I,3)
   AAA = 0.0
   DO J = 3, JMM1
      AAA = AAA + 2.0 * DELPSI / (YNEW(I,J+1) - YNEW(I,J-1))
   ENDDO
   BB = DELPSI * ( 1.0 / YNEW(I,2) + 4.0 / YNEW(I,3) )
   YOLD(I,1) = 0.4 * ( 1.0 - AAA ) / DELPSI -
   # 0.2 * ( YNEW(I,2) + YNEW(I,3) * 4.0 )
   YOLD(I,1) = 2.0 * ( 1.0 - AAA ) / BB - 1.0
   # YOLD(I,1) = 2.0 * YNEW(I,2) - YNEW(I,3) -
   # CF * (YNEW(I,3) - YNEW(I,2))
   YOLD(I,1) = (6.0*(1.0-CF)*YOLD(I,2)-(3.0-2.0*CF)*YOLD(I,3))/
   # (3.0-4.0*CF)
   WRITE(1,*) I, YOLD(I,1), YOLD(I,2)
   IF (I.EQ.IRETACH) THEN
      IF (YOLD(IRETACH,1) .LT. 0.0 .AND.
      # YOLD(IRETACH-1,1).GT.0.0) THEN
         STOP
      ENDIF
   ENDIF
   IF ( YOLD(I,1) .LT. 0.0 ) THEN
      IF (I.LT.IRETACH) THEN
         YOLD(I,1) = 0.0
      ENDIF
      IRETACH = I
   ENDIF
ENDIF
ELSE IF (I .GT. 21 .AND. J .LE. IRETACH) THEN
  IF (J .EQ. 2) THEN
    RB = 0.0
  END IF
END IF

CALCULATE THE RIGHT HAND SIDE OF THE SYSTEM EQUATION
WHICH DEPENDS ON THE INITIAL CONDITIONS (INTERIOR VALUE)

RI = -(C1 - C4) * PNEW(I-1, J)
    - (C1 + C4) * POLD(I+1, J)

# IF (J .EQ. 2) THEN
RI = RI + C2 * (PNEW(I-1, J) - POLD(I+1, J))
ELSE IF (J .EQ. JMM1) THEN
RI = RI - C2 * (PNEW(I-1, JMM1) - POLD(I+1, JMM2))
ELSE
RI = RI + C2 * (PNEW(I-1, J-1) - PNEW(I-1, J-1)
    + POLD(I+1, J-1) - POLD(I+1, J+1))
# ENDIF

IF (I .EQ. 2) THEN
RI = -(C1 + C4) * POLD(3, J)
IF (J .EQ. 2) THEN
RI = RI + C2 * POLD(3, J)
ELSE IF (J .EQ. JMM1) THEN
RI = RI + C2 * POLD(3, JMM2)
ELSE
RI = RI + C2 * (POLD(3, J-1) - POLD(3, J+1))
ENDIF
ELSE IF (I .EQ. IMM1) THEN
RI = -(C1 - C4) * PNEW(IMM2, J)
IF (J .EQ. 2) THEN
RI = RI + C2 * PNEW(IMM2, J)
ELSE IF (J .EQ. JMM1) THEN
RI = RI - C2 * PNEW(IMM2, JMM2)
ELSE
RI = RI + C2 * (PNEW(IMM2, J+1) - PNEW(IMM2, J-1))
ENDIF

ENDIF

THIS IS THE RECIRCULATING REGION AND THE REATTACHMENT POINT

IF ( .NOT. WCHECK ) THEN
  IF ( I .GT. 21 .AND. I .LT. IRETACH ) THEN
    IF ( J .EQ. 2 ) THEN
      RI = -(C1 + 2.0*C2) * PNEW(I-1, J)
      # (C1 - 2.0*C2) * POLD(I+1, J) +
      2.0 * C2 * (PNEW(I-1, J) - POLD(I+1, J))
    ENDIF
  ENDIF
C  ENDIF
C R(J) = RB + RI
100 CONTINUE
THE RECIRCULATING REGION HAS DIFFERENT MATRIX

IF (.NOT. WCHECK ) THEN
  IF ( I .GT. 21 .AND. I .LT. IRETACH ) THEN
    IF ( J .EQ. 2 ) THEN
      B(J) = -2.0 * ( C1 + C5 )
      C(J) = 2.0 * C5
    ENDIF
  ENDIF
ENDIF

CALCULATE THE RIGHT HAND SIDE OF THE SYSTEM EQUATION
WHICH DEPENDS ON THE BOUNDARY CONDITIONS (RHSi)

IF(J .EQ. 2) THEN
  RB = C2 * (PBND(I+1,1) - PBND(I-1,1))
  #
  ELSE IF(J. EQ. JMM1) THEN
  RB = C2 * (PBND(I-1,JMAX) - PBND(I+1,JMAX))
  #
  ELSE
  RB = 0.0
ENDIF

IF(I .EQ. 2) THEN
  RB = C2 * (PBND(1, J+1) - PBND(1, J-1))
  #
  ELSE IF(J. EQ. 2) THEN
  RB = RB + C2 * PBND(3,1) - (C3 - C5) * PBND(I,1)
  ENDIF
ELSE IF(I .EQ. IMM1) THEN
  RB = C2 * (PBND(IMAX, J-1) - PBND(IMAX, J+1))
  #
  ELSE IF(J. EQ. 2) THEN
  RB = RB - C2 * PBND(IM2,1)
  #
  ELSE IF(J. EQ. JMM1) THEN
  RB = RB + C2 * PBND(IM2,JMAX)
  #
  ENDIF
ENDIF

IF ( .NOT. WCHECK ) THEN

THIS IS THE SEPARATION POINT

IF ( I .EQ. 21 ) THEN
  IF(J .EQ. 2) THEN
    RB = -C2 * HSTEP
  ENDIF
ENDIF

THIS IS THE RECIRCULATING REGION AND THE REATTACHMENT POINT
CONTINUE
IF(WCHECK) THEN

CHECK THE CONVERGENCE ACCORDING TO THE OUTER CONVERGENCE CRITERIA

IF(CHECK(WNEW,WOLD,IMAX,JMAX,EPSILW)) THEN

SOLUTION HAS CONVERGED, RECORD NUMBER OF ITERATIONS AND EXIT THE DO-LOOP

N = K
GOTO 999
ENDIF
ELSE

CALCULATE THE DAMPING FACTOR FOR THE INNER CONVERGENCE CRITERIA

IF(CHECK(YNEW,YOLD,IMAX,JMAX,EPSILY)) THEN

SOLUTION HAS CONVERGED, RECORD NUMBER OF ITERATIONS AND EXIT THE DO-LOOP

N = K
GOTO 999
ENDIF

IF(K .GT. MMU) THEN
YE=0.0
DO 102 I=2,IMM1
   DO 101 J=2,JMM1
      YE=AMAX1(YE,ABS(YNEW(I,J) - YOLD(I,J)))
   101 CONTINUE
102 CONTINUE
IF(YEOLD.EQ.0.0)YEOLD=1.0
RHO = YE / YEOLD
DELTAN=RHO/(RHO+2.0)
YEOLD=YE
WRITE(1,*)(DELTAN)
ENDIF
ENDIF

TRANSFER THE NEW PHI'S TO THE OLD PHI'S

TRANSFER THE BOUNDARIES IF IT IS THE FIRST STEP

IF(WCHECK) THEN

TRANSFER THE BOUNDARIES

DO 810 I=1,IMAX
   WOLD(I,1)=WBND(I,1)
   WOLD(I,JMAX)=WBND(I,JMAX)
810 CONTINUE
**DO 710 J=1,JMAX**

  **WOLD(1,J)=WBND(1,J)**

  **WOLD(IMAX,J)=WBND(IMAX,J)**

**C710 CONTINUE**

**TRANSFER THE INTERIOR VALUES**

**DO 700 I = 2, IMM1**

  **DO 800 J = 2, JMM1**

    **WOLD(I,J) = WNEW(I,J)**

**800 CONTINUE**

**700 CONTINUE**

**CALCULATE THE RELATIVE ERROR AGAINST THE BENCH VALUES**

**WRITE(1,*),(WOLD(I,1),I=1,IMAX),ERROR**

ELSE

**TRANSFER THE BOUNDARIES**

**DO 815 I=1,IMAX**

  **YOLD(I,1)=YBND(I,1)**

  **YOLD(I,JMAX)=YBND(I,JMAX)**

**C815 CONTINUE**

**DO 715 J=1,JMAX**

  **YOLD(1,J)=YBND(1,J)**

  **YOLD(IMAX,J)=YBND(IMAX,J)**

**C**

**TRANSFER THE INTERIOR VALUES**

**DO 701 I = 2, IMM1**

  **DO 801 J = 2, JMM1**

    **YOLD(I,J) = YNEW(I,J)**

**C WRITE(*,*),I,J,YOLD(I,J)**

801 CONTINUE

701 CONTINUE

**ENDIF**

**INCREMENT NUMBER OF ITERATIONS**

**899 CONTINUE**

**N = K**

**SOLUTION FAILS TO CONVERGE**

**WRITE(1,'(A)') SOLUTION APPARENTLY FAILS TO CONVERGE**

**GOTO 1100**

**999 CONTINUE**

**W VALUES HAVE CONVERGED, PRINT SOLUTION**

**WRITE(1,199)N**

**FORMAT(' NUMBER OF ITERATIONS = ',I4)**
WRITE(1,97)
FORMAT(10X,'SOLUTION'/3(9X,'I',3X,'X',6X,'Y',2X))
DO 600 I = 1,20
   WRITE(1,99)I,X(I),YOLD(I,1),I+20,X(I+20),
   YOLD(I+20,1),I+40,X(I+40),YOLD(I+40,1)
#600 CONTINUE
99 FORMAT(3(I10,F5.1,F8.5))
1100 CONTINUE
CLOSE(UNIT=1)
STOP
END

LOGICAL FUNCTION CHECK(A,B,IL,IU,E)

   THIS FUNCTION CHECKS TO SEE IF (A(I,J) - B(I,J)) < E FOR ALL I,J

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
DIMENSION A(IL,IU), B(IL,IU)
CHECK = .TRUE.
DO 200 I = 2, IL-1
   DO 100 J = 2, IU-1
      IF(ABS(A(I,J) - B(I,J)) .GE. E) THEN
         CHECK = .FALSE.
         RETURN
      ENDIF
   100 CONTINUE
200 CONTINUE
RETURN
END

SUBROUTINE TRID(A,B,C,D,F,NL,NU)

SCALAR TRIDIAGONAL SOLVE (THOMAS ALGORITHM)
NL AND NU ARE INDEX LIMITS, NL <= N <= NU
A,B,C ARE TRIDIAGONAL ELEMENTS WITH B ARRAY ON THE MAIN DIAGONAL
D IS SCRATCH OR DUMMY ARRAY
F IS FHS FORCING FUNCTION, THE SOLUTION IS OVERLOADED IN F

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
DIMENSION A(1),B(1), C(1), D(1), F(1)
D(NL) = C(NL) / B(NL)
F(NL) = F(NL) / B(NL)
NLP = NL + 1
DO 10 N = NLP, NU
   Z = 1.0 / (B(N) - A(N) * D(N-1))
   D(N) = C(N) * Z
   F(N) = (F(N) - A(N) * F(N-1)) * Z
10 CONTINUE

BACK SWEEP

NUP = NU + NL
DO 20 NN = NLP, NU
   N = NUP - NN
   F(N) = F(N) - D(N) * F(N + 1)
20 CONTINUE
RETURN
END

double precision FUNCTION PBND(I, J)

   THIS FUNCTION RETURNS THE BOUNDARY VALUE OF AN ELEMENT OF
   THE
   PHI ARRAY WHICH COULD BE EITHER Y BOUNDARY VALUE OR W
   BOUNDARY VALUE

   COMMON /WCH/WCHECK
   LOGICAL WCHECK
   IF (WCHECK) THEN
      PBND = WBND(I, J)
   ELSE
      PBND = YBND(I, J)
   ENDIF

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PBND = POLD(I, J)
RETURN
END

double precision FUNCTION PNEW(I, J)

   THIS FUNCTION RETURNS THE FUNCTION VALUE AT A NEW ITERATION
   LEVEL OF
   PHI ARRAY WHICH COULD BE EITHER Y OR W

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX = 61, JMAX = 21)
DIMENSION WOLD(IMAX, JMAX), WNEW(IMAX, JMAX),
#  YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /VALUE/WOLD, WNEW/YVALUE/YOLD, YNEW/WCH/WCHECK
LOGICAL WCHECK
IF (WCHECK) THEN
   PNEW = WNEW(I, J)
ELSE
   PNEW = YNEW(I, J)
ENDIF
RETURN
END

double precision FUNCTION POLD(I, J)

   THIS FUNCTION RETURNS THE FUNCTION VALUE AT A OLD ITERATION
   LEVEL OF
   PHI ARRAY WHICH COULD BE EITHER Y OR W

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX = 61, JMAX = 21)
DIMENSION WOLD(IMAX, JMAX), WNEW(IMAX, JMAX),
# YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /VVALUE/WOLD, WNEW/YVALUE/YOLD, YNEW/WCH/WCHECK
LOGICAL WCHECK
IF (WCHECK) THEN
   POLD = WOLD(I, J)
ELSE
   POLD = YOLD(I, J)
ENDIF
RETURN
END

double precision FUNCTION YBND(I, J)

THIS FUNCTION RETURNS THE BOUNDARY VALUE OF Y AT I, J.
I MUST EQUAL 1 OR IMAX, OR J MUST EQUAL 1 OR JMAX

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX=61, JMAX = 21)
DIMENSION YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /YVALUE/YOLD, YNEW
COMMON /BND/DELX, DELPSI, XMIN, PSIMIN, REY
DATA PI/3.1415927/
IF (I .EQ. 1) THEN
   THETA = ACOS(1.0 - 2.0*PSI(J)) / 3.0
   YBND = COS(THETA + 4.0 * PI / 3.0) + 0.9
ELSE IF (I.EQ.IMAX) THEN
   THETA = ACOS(1.0 - 2.0*PSI(J)) / 3.0
   YBND = 1.4 * COS(THETA + 4.0 * PI / 3.0) + 0.7
ELSIF (J .EQ. 1) THEN
   IF ( I .LE. 21 ) THEN
      YBND = 1.4
   ELSE IF ( I .GE. 29 ) THEN
      YBND = 0.0
   ELSE
      YBND = 0.4 * SQRT( 1.0 - ( (X(I) - 2.0) / 0.8 ) ** 2 )
   ENDIF
ENDIF
RETURN
END

double precision FUNCTION WBND(I, J)

THIS FUNCTION RETURNS THE BOUNDARY VALUE OF W AT I, J.
I MUST EQUAL 1 OR IMAX, OR J MUST EQUAL 1 OR JMAX

IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX=61, JMAX = 21)
DIMENSION WOLD(IMAX, JMAX), WNEW(IMAX, JMAX),
# YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /VVALUE/WOLD, WNEW/YVALUE/YOLD, YNEW/WA/HSTEP
IF (I .EQ. 1) THEN
    WBN = 12.0 * YBN(I,J) - 10.8
ELSE IF (I.EQ.IMAX) THEN
    WBN = WOLD(I-1,J)
C4.37177843 * YBN(I,J) - 3.06122449
ELSE IF (J.EQ.1) THEN
    IF ( I .GT. 21 .AND. I .LT. 29 ) THEN
        WBN = 0.0
    ELSE
        WBN = 3.0 * ( FLOAT(I-1) / 60.0 - 2.0 )
    ENDIF
ELSE IF (J.EQ.JMAX) THEN
    WBN = WOLD(I,J-1)
C-3.0 * ( FLOAT(I-1) / 60.0 - 2.0 )
ENDIF
RETURN
END

double precision FUNCTION PSI(J)
IMPLICIT DOUBLE PRECISION (A-H, O-Z)
COMMON /BND/DLX, DELPSI, XMIN, PSIMIN, REY
PSI = PSIMIN + (J - 1) * DELPSI
RETURN
END

double precision FUNCTION X(I)
IMPLICIT DOUBLE PRECISION (A-H, O-Z)
COMMON /BND/DLX, DELPSI, XMIN, PSIMIN, REY
X = XMIN + (I - 1) * DLX
RETURN
END

double precision FUNCTION DELXY(I,J)
IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX=61, JMAX=21)
DIMENSION YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /BND/DLX, DELPSI, XMIN, PSIMIN, REY
COMMON /YVALUE/YOLD, YNEW
IF (I .EQ. 1) THEN
    DELXY = (YOLD(I+1, J) - YOLD(I, J)) / DLX
ELSE IF (I .EQ. IMAX) THEN
    DELXY = (YOLD(I, J) - YOLD(I-1, J)) / DLX
ELSE
    DELXY = (YOLD(I+1, J) - YOLD(I-1, J)) / (2.0 * DLX)
ENDIF
RETURN
END

double precision FUNCTION DELPSY(I,J)
IMPLICIT DOUBLE PRECISION (A-H, O-Z)
PARAMETER (IMAX=61, JMAX=21)
DIMENSION YOLD(IMAX, JMAX), YNEW(IMAX, JMAX)
COMMON /BND/DLX, DELPSI, XMIN, PSIMIN, REY
COMMON /YVALUE/YOLD, YNEW
IF (J .EQ. 1) THEN
  DELPSY = (YOLD(I, J+1) - YOLD(I, J)) / DELPSI
ELSE IF (J .EQ. JMAX) THEN
  DELPSY = (YOLD(I, J) - YOLD(I, J-1)) / DELPSI
ELSE
  DELPSY = (YOLD(I, J+1) - YOLD(I, J-1)) / (2.0 * DELPSI)
END IF
RETURN
END
VITA AUCTORIS

Paul Carson graduated from the University of Windsor in Honours Mathematics in 1969 (B.Sc.) and in Applied Mathematics in 1970 (M.Sc.). He received a Master's Degree in Aeronautical Engineering (M.Eng.) in 1977 from Carleton University and his P.Eng. from the Association of Professional Engineers in 1981. He holds both senior Airline Transport Pilot Licences (ATPL) issued by Canada and the United States. He has worked for the Defence Research Board of Canada, the Federal Department of Transport and Bell Canada in a number of capacities, including research scientist, engineer and pilot. Mr. Carson is a member of several professional organizations and associations including the Canadian Aeronautics and Space Institute, the Association of Professional Engineers of Ontario, the Canadian Airline Pilots' Association, the International Society for Air Safety Investigators, and most recently, the Human Factors Association of Canada. He is currently studying towards an honours B.A. degree in Psychology at the University of Ottawa.