The vortex-entrainment sheet in an inviscid fluid: theory and separation at a sharp edge

A. C. DeVoria$^1$ and K. Mohseni$^{1,2}$

$^1$Department of Mechanical & Aerospace Engineering, University of Florida, Gainesville, FL 32611, USA
$^2$Department of Electrical & Computer Engineering, University of Florida, Gainesville, FL 32611, USA

(Received xx; revised xx; accepted xx)

In this paper a model for viscous boundary and shear layers in three-dimensions is introduced and termed a vortex-entrainment sheet. The vorticity in the layer is accounted for by a conventional vortex sheet. The mass and momentum in the layer are represented by a two dimensional surface having its own internal tangential flow. Namely, the sheet has a mass density per-unit-area making it dynamically distinct from the surrounding outer fluid and allows the sheet to support a pressure jump. The mechanism of entrainment is represented by a discontinuity in the normal component of velocity across the sheet. The velocity field induced by the vortex-entrainment sheet is given by a generalized Birkhoff-Rott equation with a complex sheet strength. The model was applied to the case of separation at a sharp edge. No supplementary Kutta condition in the form of a singularity removal is required as the flow remains bounded through an appropriate balance of normal momentum with the pressure jump across the sheet. A pressure jump at the edge results in the generation of new vorticity. The shedding angle is dictated by the normal impulse of the intrinsic flow inside the bound sheets as they merge to form the free sheet. When there is zero entrainment everywhere the model reduces to the conventional vortex sheet with no mass. Consequently, the pressure jump must be zero and the shedding angle must be tangential so that the sheet simply convects off the wedge face. Lastly, the vortex-entrainment sheet model was demonstrated on two shedding example problems.

1. Introduction

An inviscid fluid is governed by the Euler equations, which permit surfaces of discontinuity or jumps in physical quantities as part of the solution. While these solutions are only mathematical idealizations, their study is of practical use because they can serve as approximations to physical viscous phenomena, particularly at high-Reynolds number ($Re$). For example, combinations of jumps in the fluid density, velocity, pressure, entropy can be cast to represent shocks, boundary/shear layers, or interfaces between fluid media. Here, we focus on the inviscid modeling of viscous layers and flow separation, which has long been a challenging interest in the general fluid dynamics community.

Typically, the high-$Re$ or inviscid approximation $Re \to \infty$ of viscous layers has been the infinitely thin vortex sheet. Namely, the actual distribution of the velocity and vorticity fields within the layer are lost and the sheet is characterized by a jump in the tangential velocity, which is the local sheet strength representing the integral of vorticity across

† Email address for correspondence: mohseni@ufl.edu
the layer. In this way, the vortex sheet strength preserves the total circulation. Stated differently, the vortex sheet strength is the circulation per-unit-length of sheet.

In general, the evolution equation for a conventional vortex sheet is given by a Biot-Savart integral. Pozrikidis (2000) discusses several of the difficult aspects of computing three-dimensional vortex sheet dynamics including treatments of the principal value integral. For a two-dimensional domain this integral is known as the Birkhoff-Rott equation when expressed in a complex analysis formulation. Pullin (1978) used this approach and obtained a similarity solution for the shape of a rolled-up semi-infinite free vortex sheet as well as the shapes of vortex sheets shed from the apex of infinite wedges. Baker et al. (1982) used a vortex sheet to represent the interface between a stratified flow to study breaking water waves. Jones (2003) computed the truly unsteady vortex sheet shedding from the cusped edges of a moving, finite-chord flat plate where the plate was also represented as a vortex sheet of known geometry. DeVoria & Mohseni (2018) proposed a model of the local self-induced velocity of a vortex sheet segment that allows integration through the singularity of the Birkhoff-Rott equation.

Calculation of the dynamics of vortex sheets may be simplified with an approximation obtained by discretizing the vortex sheet into point vortices, a technique that dates back to Rosenhead (1931) and Westwater (1935) in the 1930s and represents the beginning of what is now referred to as vortex methods (Leonard 1980). Multi-vortex shedding models have become commonplace (e.g., Clements 1973; Sarpkaya 1975; Katz 1981; Cortelezzi & Leonard 1993; Nitsche & Krasny 1994; Haroldsen & Meiron 1998; Michelin & Llewellyn Smith 2009; Xia & Mohseni 2013; Wang & Eldredge 2013). These models have been used to great effect for fairly accurate predictions in physical problems, especially the load prediction on and shedding of vorticity from airfoils.

For the conventional vortex sheet the statement of mass conservation is given by the boundary condition of a continuous normal velocity (Saffman 1992), namely a no-flux condition. Therefore, vortex sheets are mass-less ‘contact discontinuities’ (Wu et al. 2006) that remain dynamically indistinct from the surrounding fluid(s). However, real boundary/shear layers contain mass and momentum, which are entrained into the layer. Here, we propose a model that provides a fuller description of the dynamics of a viscous layer by explicitly representing the mass and momentum within the layer. Since entrainment is an inherently viscous process, such a model relaxes the constraint of $Re \to \infty$ and could be applicable to a range of Reynolds numbers provided the entrainment is modeled sufficiently well.

A physical example that highlights the need for a dynamic model is given by the shedding angle at the non-cusped trailing edge of an airfoil. Basu & Hancock (1978) considered the shedding of a conventional vortex sheet and, based on the works of Giesing (1969) and Maskell (1971), concluded that the shedding angle must be tangential to one of the edge faces. This result was also proven by Pullin (1978) for the self-similar shedding problem. On the other hand, physical intuition suggests that the shedding angle ought to vary continuously between the tangential limits, which was experimentally observed by Poling & Telionis (1986). In this paper we show that if the sheet contains mass then a non-tangential shedding angle is indeed possible.

Xia & Mohseni (2017) recently studied the formation of vortex sheets at the trailing edge of airfoils. They applied conservation laws to a wye-shaped control volume encompassing the two boundary layers at the trailing edge that merge to create the free shear layer. In addition to a conventional vortex sheet with a jump in tangential velocity, they also proposed a generalized sheet model for viscous layers, whereby sheets that have jumps in other physical quantities, for example a jump in the stream function representing entrainment into the layer, can be superimposed onto each other.
The objective of this study is to present and formalize the theory for the surface of discontinuity that we term the vortex-entrainment sheet. In regard to flow separation, the focus in this paper will be on separation at a sharp edge. The paper is organized as follows. In §2 the vortex-entrainment sheet is defined in a three-dimensional domain as a dividing surface containing mass and momentum and thus has its own intrinsic flow within the sheet. §§2.1-2.3 present the dynamics and kinematics of the sheet as well as the coupling of the internal flow with the flow outside the sheet. The model is then applied as a representation of the self-similar Falkner-Skan boundary layers in §3. Followed by this are two complementary methods of solving the outer flow via a Laplace equation in §4.1 and a boundary integral formulation in §4.2. A general solution algorithm is discussed in §4.3. The conditions that determine shedding/separation of a sheet from a sharp edge are given in §5. Lastly, some example calculations are presented in §6.

2. The vortex-entrainment sheet

Here, we propose the vortex-entrainment sheet as a dynamic inviscid model of viscous layers. To exemplify the physics of entrainment consider the viscous wake left behind a body, where the drag on the body may be determined from a wake survey (Batchelor 1967, pp. 348-353). The expression for the (steady) drag is proportional to the net entrainment rate from infinity, say \( Q_\infty \), into the wake: \( D = \rho U Q_\infty \). Paired with a similar expression for the lift \( L = \rho U \Gamma \), where \( \Gamma \) is the circulation, these two conjugate relations have been referred to as the Joukowski-Filon formulae (Liu et al. 2015). There are many explanations of d’Alembert’s paradox, but this drag law offers a clear interpretation of resistive forces as due to entrainment of fluid owing to the action of viscosity.

Hence, \( Q_\infty \) implicitly depends on the Reynolds number, \( Re \), through the viscous entrainment process. For example, the wake of a flat plate at zero incidence was calculated by Goldstein via the boundary layer method (Schlichting 1955, pp. 138-142). The total drag on the plate of length \( l \) and width \( w \) is \( D = \rho U (1.33 U w \sqrt{\nu / U}) \) from which we identify and rewrite \( Q_\infty = 1.33 w \nu \sqrt{Re_t} \) where \( Re_t = U l / \nu \). The explicit appearance of the viscosity \( \nu \) reiterates that entrainment is an inherently viscous process and, in general, we ought to expect that the boundary conditions will be functions of \( \nu \) (see Saffman 1981, p. 53). Hence, the drag could be recovered by equating the net entrainment rate of the vortex-entrainment sheet representing the wake to the value \( Q_\infty \). This overly simple example conceptually demonstrates how the vortex-entrainment sheet could be used to provide an improved model of viscous layers.

As an initial step towards defining the vortex-entrainment sheet we consider the mass within a control volume surrounding a viscous layer and seek to preserve the mass within
the layer as its thickness $h$ is collapsed to zero; figure 1. This concept was applied by Xia & Mohseni (2017) to the case of an airfoil trailing edge that resulted in a discontinuity of the stream function, or equivalently the normal velocity, which represents entrainment into the sheet. The distribution of fluid with density $\rho$ across the layer is preserved by assigning a surface mass density $\rho_s$ with units of mass per-unit-area of sheet. The actual distribution of flow $u$ in the layer becomes an ‘intrinsic’ flow $v$ confined on the sheet and may possess properties that are different from its bulk surroundings. Surfaces of this type commonly appear in other fields, such as electrodynamics (e.g. a capacitor with surface charge density and discontinuity in the normal component of the electric field, Jackson 1998). While such surfaces have also been considered in fluid mechanics, for example the Boussinesq-Scriven surface fluid model (Scriven 1960), they are less frequently encountered.

We will assume that the flow outside the sheet $u$ is an irrotational, incompressible flow that may be discontinuous across the sheet. However, we note that this is not to be interpreted as a $Re \rightarrow \infty$ limit, but rather that the viscous, rotational portion of the fluid has been ‘cut out’ and sutured up by the vortex-entrainment sheet. In this regard, it is especially important to note that the vortex-entrainment sheet is not a streamline. Next, in §§2.1-2.2 we more rigorously define the vortex-entrainment sheet.

2.1. Sheet dynamics

Let $x$ be the position vector of an arbitrary point in three-dimensional space. The sheet is immersed in an irrotational, incompressible fluid with density $\rho$ and velocity $u(x, t)$, where $t$ is time. A right-handed orthogonal coordinate system $(s, n, b)$ is defined on the sheet such that the basis $\hat{s}$-$\hat{b}$ spans the local tangent manifold that is the sheet and $\hat{n}$ is the corresponding sheet normal vector. The sheet is a two-dimensional surface whose location is specified by $x = x_s(s, b, t)$; see figure 2.

The jump in a quantity $f$ across the sheet is defined as

$$[f] = f^+ - f^-,$$

where $f^+$ and $f^-$ are the limiting values of $f$ as approached from the different sides of the sheet with $\hat{n}$ pointing to the (+) side. The surface velocity on the sheet is $v(x_s, t)$ and given our choice of coordinates this can be decomposed as

$$v = w + (v \cdot \hat{n})\hat{n} = (v_s \hat{s} + v_b \hat{b}) + v_n \hat{n}.$$
where \( w(x_s, t) = v_s \hat{s} + v_b \hat{b} \) is the intrinsic velocity of the flow inside the sheet and \( v_n(x_s, t) = v \cdot \hat{n} \) is the normal velocity component. The surface operator \( \nabla_s \), which only acts in the tangent space of the sheet, is given by:

\[
\nabla_s = \left( \frac{1}{h_s} \frac{\partial}{\partial s} \right) \hat{s} + \left( \frac{1}{h_b} \frac{\partial}{\partial b} \right) \hat{b},
\]

(2.3)

with \( h_s \) and \( h_b \) as the scale factors for the chosen surface coordinate system. The operator giving the surface material derivative of the quantity \( f \) on the sheet is defined as:

\[
\frac{D_s f}{Dt} = \frac{\partial f}{\partial t} + w \cdot \nabla_s f.
\]

(2.4)

Given our specific choice of the coordinate system we have \( (v_n \hat{n}) \cdot \nabla_s f = 0 \) and thus we may write \( w \cdot \nabla_s f = v \cdot \nabla_s f \). The reader is referred to Aris (1962) and Slattery et al. (2007) for further details regarding these surface operators as well as the sheet conservation equations discussed next.

Now, we define the mass within the sheet by integrating across the layer in the normal direction to give

\[
\rho_s(x_s, t) = \int \rho(x, t)dn.
\]

(2.5)

The quantity \( \rho_s \) is the sheet mass density and has units of mass per-unit-area of sheet. In this way, the mass that was originally in the viscous layer is preserved by ‘collapsing’ it to the sheet; recall figure 1. Fluid may enter the sheet via entrainment and the mass conservation equation on the sheet is:

\[
\frac{D_s \rho_s}{Dt} + \rho_s(\nabla_s \cdot v) = -[\rho(u - v) \cdot \hat{n}] .
\]

(2.6)

The left side has the familiar form of mass conservation in the bulk fluid, but with the understanding of the intrinsic sheet-tangent operators. In contrast, the right side is non-zero with \(-[\rho(u - v) \cdot \hat{n}]\) representing a source of \( \rho_s \) due to entrainment from the outer flow. While the fluid velocity may have a discontinuity, the sheet velocity is continuous meaning that \([v] = 0\).

The conservation of momentum on the sheet also takes on a similar form to the usual equation for the bulk fluid, but again with additional source terms associated with the flux and entrainment of momentum from the flow outside the sheet:

\[
\rho_s \frac{D_s v}{Dt} - (\nabla_s \cdot T_s + \rho_s f_s) = -[\rho(u - v)(u - v) \cdot \hat{n}] - [p]\hat{n} + [\tau],
\]

(2.7)

where \([p]\) is the pressure jump across the sheet and \( \tau = [T_v \cdot \hat{n}] \) is the shear stress vector with \( T_v \) as the deviatoric stress tensor outside the sheet. In general, the sheet may also have its own phenomena, such as a surface stress tensor \( T_s \) (yielding a surface tension force, for example), or a ‘body’ force \( f_s \) that acts only on the sheet mass density \( \rho_s \), as well as different surface fluid properties (Slattery et al. 2007). For simplicity we shall neglect each of these terms in (2.7). The surface flow \( v \) is clearly dynamically coupled to the outer flow \( u \), which must also be solved. The main unknowns to be solved for are \( u, \rho_s, v \) and \([p]\). For a free sheet we have \( \tau = \tau^- = 0 \) since the outer flow is potential. When the sheet is bound to a surface its normal velocity will be known from the motion \( U \) of that surface as \( v_n = \hat{n} \cdot U \) and the wall shear stress \( \tau_w = [\tau] \) is then appended to the list of unknowns.

At this stage it is useful to discuss the occurrence and interpretation of entrainment by examining the jump source term on the right side of (2.6). Using the properties of
jumps and averages of products (e.g. Wu et al. 2006, p. 34) we may write
\[ \rho(u_n - v_n) = \rho(\overline{u_n} - v_n) + \rho[u_n], \]  
(2.8)
where \( \overline{\cdot} = [(\cdot)^+ + (\cdot)^-]/2 \) is the arithmetic mean of the values of a quantity across the sheet. Throughout, an overline will indicate this type of average unless otherwise noted. The sheet must coincide with the discontinuity \( [u_n] \neq 0 \) for entrainment to occur meaning that fluid will not pass through the sheet. However, since fluid may be entrained into the sheet, then \( v_n \) will take a value between \( u_n^+ \) and \( u_n^- \). To elucidate this, consider a weighted average:
\[ v_n = \left[ \frac{1+\alpha}{2} \right] u_n^+ + \left[ \frac{1-\alpha}{2} \right] u_n^- = \overline{u_n} + \frac{\alpha}{2} [u_n], \]
(2.9)
where \( \alpha \) is a weighting parameter with \(-1 \leq \alpha \leq 1\) as similarly used by Baker et al. (1982). Substituting (2.9) into the preceding equation gives:
\[ [\rho(u_n - v_n)] = [u_n] (\overline{p} - \frac{\alpha}{2} [\rho]) = [u_n] ([\frac{1-\alpha}{2}] \rho^+ + [\frac{1+\alpha}{2}] \rho^-) \]
(2.10)
so that fluid with a weighted average density is entrained at a rate \( [u_n] \). However, notice the weighting is opposite to that of \( v_n \); for example, if \( \alpha > 0 \) then \( v_n \) moves with a velocity closer to \( u_n^+ \), but in doing so the sheet entrains or ‘consumes’ more fluid with the density \( \rho^- \). In a similar fashion, substitution of (2.9) into the momentum flux term in (2.7) yields:
\[ [\rho(u - v)(u_n - v_n)] = [u_n] \left( \rho(u - v) - \frac{\alpha}{2} [\rho(u - v)] \right) \]
(2.11)
so that \( [u_n] \) simultaneously entrains a weighted average of the difference momentum \( \rho(u - v) \) from either side of the sheet. If there is no mass in a free sheet, \( \rho_s = 0 \), then (2.6) dictates that \( [u_n] = 0 \). Then (2.7) gives a zero pressure jump \( [p] = 0 \). Of course, these are the kinematical and dynamical features of a conventional vortex sheet. Note that each was consequently obtained from the assumption that \( \rho_s = 0 \).

2.2. Sheet kinematics

In order to solve the surface flow equations in (2.6)-(2.7) we must also determine the outer flow \( u \). A vector field is determined from its essential characteristics, namely its divergence, curl, discontinuities and boundary values (Phillips 1959). In general, the field can be expressed by the sum of a divergence-free part and a curl-free part via a Helmholtz-Hodge decomposition, and if the full vector is known on the boundary a harmonic part can be distinguished (Bhatia et al. 2013). As such, the decomposition can be made unique in several ways depending on the known boundary conditions.

While \( u \) is an irrotational, incompressible flow with no distribution of vorticity (\( \omega \)) or dilatation rate (\( \Delta \)) by definition, it is discontinuous across the sheet. Therefore, the decomposition of \( u \) has distinct contributions from its discontinuities. These contributions can be rigorously derived (e.g. Phillips 1959), but for simplicity we employ a straightforward, albeit crude device that yields the same result. If we insist on allowing the fluid to possess a singular curl and divergence we can represent the discontinuity in \( u \) across the sheet with a Heaviside function. Then the singular parts of the curl and divergence are (see Appendix A):
\[ \omega = \nabla \times u = (\hat{n} \times [u])\delta(n), \quad \Delta = \nabla \cdot u = (\hat{n} \cdot [u])\delta(n), \]
(2.12)
where \( \delta(n) \) is the Dirac delta function and \( n \) is the sheet normal coordinate. These components of the curl and contribution to the divergence are those affected by normal derivatives across the sheet. The intrinsic surface operators ‘remove’ these singular parts.
respectively, with effectively take the curl and divergence of (2.17) by the cross product and dot product, and we have used the identity \( \mathbf{\nabla} \cdot \mathbf{u} = \frac{1}{2} \left| \mathbf{u} \right|^2 + \frac{P}{\rho} \), since normal derivatives across the sheet are not defined. Therefore, we isolate the tangential and normal components of the right side by making use of the surface-tangent operator \( \mathbf{\nabla}_s \) in (2.3) to allow the jump operator inside on the tangential part:

\[
\left[ \mathbf{u} \cdot \mathbf{\nabla} \mathbf{u} + \mathbf{\nabla}_s \frac{P}{\rho} \right] = \mathbf{\nabla}_s \left[ \frac{1}{2} \left| \mathbf{u} \right|^2 + \frac{P}{\rho} \right] + \left[ \mathbf{\nabla} \cdot (u_n \mathbf{u}) + \mathbf{\hat{n}} \cdot \mathbf{\nabla}_s \frac{P}{\rho} \right] \mathbf{\hat{n}},
\]

and we have used the identity \( \mathbf{u} \cdot \mathbf{\nabla} \mathbf{u} = \mathbf{\nabla} \cdot (\mathbf{u} \otimes \mathbf{u}) \) for the normal component. We can effectively take the curl and divergence of (2.17) by the cross product and dot product, respectively, with \( \mathbf{\hat{n}} \). Using the sheet strength definitions in (2.13)-(2.14) along with the
properties of jumps of product quantities and the triple vector product we have the following identities:

\[ \left[ \frac{1}{2} |\mathbf{u}|^2 \right] = \mathbf{u} \cdot [\mathbf{u}] = \mathbf{u} \cdot (\gamma \times \hat{\mathbf{n}} - q\hat{\mathbf{n}}). \]  

(2.19)

We then obtain:

\[ \hat{\mathbf{n}} \times \left[ \mathbf{u} \cdot \nabla \mathbf{u} + \nabla p_\rho \right] = \frac{\partial}{\partial b} \left( \mathbf{u} \cdot [\mathbf{u}] + \left[ \frac{p}{\rho} \right] \right) \hat{s} - \frac{\partial}{\partial s} \left( \mathbf{u} \cdot [\mathbf{u}] + \left[ \frac{p}{\rho} \right] \right) \hat{b} \]  

(2.20)

\[ \hat{\mathbf{n}} \cdot \left[ \mathbf{u} \cdot \nabla \mathbf{u} + \nabla p_\rho \right] = \frac{\partial}{\partial s} \left( u_n u_s \right) + \frac{\partial}{\partial b} \left( u_n u_b \right) + \left[ \frac{\partial}{\partial n} \left( u_n^2 + \frac{p}{\rho} \right) \right]. \]  

(2.21)

In the event that there is no entrainment, then (2.21) would be zero since the pressure jump immediately communicates normal momentum from one side of the vortex-entrainment sheet to the other. To see this, consider the normal component of the sheet momentum equation (2.7):

\[ \rho_s \frac{D_s v_n}{Dt} = -\left[ \rho (u_n - v_n)^2 \right] - [p] = -2\rho (u_n - v_n) [u_n] - [p]. \]  

(2.22)

With \( q = -[u_n] = 0 \) then \([p]\) delivers the acceleration \( D_s v_n / Dt \) to the sheet mass \( \rho_s \). On the other hand, when \( q \neq 0 \) then the last term in (2.21) represents a gradient of energy across the sheet meaning that the pressure jump must expend additional energy to also accelerate the newly entrained sheet mass. Obviously we cannot evaluate this gradient directly. However, from what has been said about the \( q = 0 \) case we can conclude that the normal component of the jump in the Euler equation is equal to the amount of normal momentum that is lost or entrained into the sheet. Therefore, after defining

\[ \mu \equiv \mathbf{u} \cdot [\mathbf{u}] + \left[ \frac{p}{\rho} \right] = \mathbf{u} \cdot (\gamma \times \hat{\mathbf{n}} - q\hat{\mathbf{n}}) + \left[ \frac{p}{\rho} \right]. \]  

(2.23)

we recombine the tangential and normal components to obtain the evolution equation for the vectorial strength \( \alpha = \gamma + q\hat{\mathbf{n}} \) of the vortex-entrainment sheet as:

\[ \rho_s \frac{\partial \alpha}{\partial t} = (\rho_s \nabla_s \mu) \times \hat{\mathbf{n}} + \left( \left[ \rho (u_n - v_n)^2 \right] \right) \hat{n}. \]  

(2.24)

Evidently the evolution of the sheets are coupled to each other through the quantity \( \mu \), which partly represents the jump in dynamic pressure. For the conventional vortex sheet (2.24) represents the familiar result that tangential pressure gradients generate vorticity components bi-normal to the gradient direction (Lighthill 1963). For the entrainment sheet the pressure jump acts to ‘push’ or entrain fluid into the sheet.

2.4. Physics of inviscid separation: interpretation of the Kutta condition

We now briefly discuss inviscid flow separation at a sharp edge from a physical point of view. Regardless of the shedding problem or its formulation, the requirement of a bounded flow at the sharp edge should be imposed. This constraint is the Kutta condition, which we note has no fundamental basis. However, from (2.21)-(2.22) we see that the condition of a regular flow is ensured by an appropriate balance of normal momentum with the pressure jump. For the conventional vortex sheet with \([u_n] = [p] = 0\) this is automatically satisfied by the condition of mass conservation, namely \( \rho_s = 0 \).

Next, consider the inviscid mechanism that eliminates the singular pressure \textit{gradient}, but not the jump, and allows the flow to separate from the surface. To see how this neutralization of the normal pressure gradient to its jump value relates to vorticity, we again employ the simple device of allowing the fluid to have a singularity in the form of
a delta function as was done for the curl and divergence in (2.12). Namely, the singular part of the pressure gradient is $\nabla (p/\rho) = \mathbf{n}[p/\rho] \delta(n)$. Since $p$ is not $C^2$ at the irregular sharp edge point, then $\nabla \times \nabla (p/\rho)$ is not necessarily zero there (Greenberg 1998) and thus contributes to the (singular) vorticity equation for $\omega = \gamma \delta(n)$ as:

$$\nabla \times \nabla \frac{p}{\rho} = \left( \frac{\partial}{\partial b} \left[ \frac{p}{\rho} \right] \hat{s} - \frac{\partial}{\partial s} \left[ \frac{p}{\rho} \right] \hat{b} \right) \delta(n),$$

(2.25)

which are the pressure terms on the right side of (2.20). Hence, a non-zero pressure jump at the edge generates new vorticity and will result in a loading there (e.g. see Sears 1956). For the conventional vortex sheet $[p] = 0$ and while $\omega$ is still singular the above pressure terms are in fact zero. Therefore, in this case the vorticity in a vortex sheet simply convections off the edge.

Lastly, an interesting physical interpretation of separation arises when a wall-bounded vortex-entrainment sheet is represented as a dipole sheet or ‘double layer’ (Kellogg 1929). The singularity in the pressure gradient field at the edge can be represented by adding the potential of a doublet to the vector potential decomposing $\nabla p$ (Phillips 1959). Hence, the finite strength of the doublet would be equal to the pressure jump. Then the neutralization of the pressure gradient singularity ‘tears apart’ the double layer and sheds one sheet into the fluid and the other sheet inside the wall as its image.

3. Falkner-Skan boundary layers

Before moving on to the solution of the outer flow in §4, we apply the vortex-entrainment sheet model to the self-similar Falkner-Skan boundary layer solutions. This also serves the purpose of revealing details of an actual viscous entrainment process. We will obtain the sheet strengths $\gamma$ and $q$ directly from the solutions. The driving outer flow $U(x) = ax^m$ is from left-to-right and the local Reynolds number is $Re_x = Ux/\nu$ and is based on the distance $x$ from the leading edge.

On the solid surface we have $u = u_s^- = 0$ and $v = u_n^- = 0$ due to the no-slip and no-penetration conditions. The Cartesian components of velocity at the outer edge of the boundary layer are $u^+ = U(x)$ and $v^+ = c_x U(x)/\sqrt{Re_x}$ where $c_x$ is a numerical coefficient. In order to calculate the proper vortex and entrainment sheet strengths given in (2.13)-(2.14), we must use the tangential and normal components. In other words, we use the known shape of the boundary layer, which we write as $\delta(x) = c_b x/\sqrt{Re_x}$. Hence, $u_s^+ = u^+ \cdot \hat{s}^+$ and $u_n^+ = u^+ \cdot \hat{n}^+$ where $\hat{s}^+ = (\cos \theta, \sin \theta)$ and $\hat{n}^+ = (-\sin \theta, \cos \theta)$ with $\tan \theta = d\delta/dx$. After some algebra we obtain:

$$\gamma(x) = 0 - u_s^+ = -U(x) \left[ \frac{\sqrt{Re_x} + c_m \frac{c_x c_b}{\sqrt{Re_x}}}{\sqrt{(c_b c_m)^2 + Re_x}} \right],$$

(3.1)

$$q(x) = 0 - u_n^+ = U(x) \left[ \frac{c_b c_m - c_x}{\sqrt{(c_b c_m)^2 + Re_x}} \right],$$

(3.2)

where $c_m = (1-m)/2$. When $Re_x \to \infty$ the entrainment dies out and the vortex sheet strength becomes $-U(x)$ as expected from the infinite-Reynolds number assumption associated with vortex sheets. As $Re_x$ decreases the entrainment strength becomes relatively more significant. However, the limit $Re_x \to 0$ cannot be interpreted with any real physical meaning since the boundary layer equations breakdown and are invalid at the leading edge, which leads to the singular behavior of the vortex sheet strength.

Hence, we can safely assume the region of validity to begin at some downstream location
such that $Re_x \gg (c_3c_m)^2$. We computed numerical solutions of the Falkner-Skan equation for $-0.0904 \leq m \leq 2$ and indeed find that $|c_3c_m| \sim o(10)$; the edge of the boundary layer was defined with the 99% rule. Using this approximation and upon substituting the definitions of the coefficients $c_3$ and $c_v$ the $\sqrt{Re_x}$ factor cancels out and the sheet strengths can be written in terms of the non-dimensional boundary layer thickness $\delta/x$ and vertical velocity $v^+/U$ at the edge of the layer as:

$$
\gamma(x) \approx -U(x) \left[ 1 + c_m \frac{\delta}{x} \frac{v^+}{U} \right], \quad q(x) \approx U(x) \left[ c_m \frac{\delta}{x} - \frac{v^+}{U} \right].
$$

The first and second terms in the square brackets of each expression correspond to the horizontal and vertical Cartesian components $u^+=U(x)$ and $v^+(x)$, respectively. For $q$ there is a competition between the acceleration of the imposed flow ‘pushing’ fluid into the boundary layer and the growth of the layer displacing fluid upwardly.

Figure 3(a) plots $c_3$ and $c_v$ against $m$ along with $c_q \equiv c_m c_3 - c_v$, which is the numerator of (3.2) and represents the sign of $q$, which is positive for any $m$. The decreasing trend in $c_3$ is the collapsing boundary layer thickness corresponding to a stronger outer flow and thus larger $Re_x$. For decelerating and slowly accelerating flows $c_v > 0$ meaning that displacement outpaces entrainment. There is a positive vertical velocity at the edge of boundary layer and thus an efflux of fluid at infinity. For larger accelerations the trend reverses with $c_v, v^+ < 0$ above $m \approx 0.2$ and there is an influx flow at infinity to compensate for the entrainment.

Next, we consider the intrinsic flow in the vortex-entrainment sheet representing a boundary layer. For simplicity we take the Blasius boundary layer with $m = 0$ and so $U(x) = U_o$. Noting that the problem is two-dimensional, steady and $v_n = 0$ since the sheet is bound to the stationary plate ($u^+_s = u^+_n = 0$), then the mass and momentum equations in (2.6)-(2.7) reduce to:

$$
\frac{\partial}{\partial s} (\rho_s u_s) = \rho q, \quad \tau_w = \rho q u_s + \rho \gamma u_n - \frac{\partial}{\partial s} (\rho_s v^+_s)^2, \quad [p] = 2\rho u_n q.
$$

Note that the sheet coordinate is also the plate coordinate $s = x$. The sheet strengths $\gamma$ and $q$ are obtained from (3.1)-(3.2) and we can use the plate boundary conditions to give $\overline{u_s} = u^+_s/2 = -\gamma/2$ and $\overline{u_n} = u^+_n/2 = -q/2$. Since the entrainment strength decays significantly for $x \gg 1$, the normal momentum equation gives $[p] \approx 0$ which is
the boundary layer assumption that the pressure is constant across the layer, or more precisely $O(Re_x^{-1/2})$. Of course this small pressure jump due to entrainment is balanced by an equal and opposite one on the other side of the plate.

For a time-dependent problem, we must be given an initial condition for $\rho_s$ and $v_s$. In this steady case we can obtain the distribution of $\rho_s$ from (2.5) so that $\rho_s = \rho \delta$. Integrating the mass equation yields $\rho_s(x)v_s(x)$ and subsequently we obtain:

$$v_s(x) - v_s(0) = \frac{1}{\rho_s(x)} \int_0^x p q ds = \frac{1}{c_s \sqrt{Re_x}} \int_0^{Re_x} q d(Re_s)$$

$$= \frac{2U_0(c_s c_m - c_v)}{c_s \sqrt{Re_x}} \left[ \sqrt{(c_s c_m)^2 + Re_x} - Re_m c_m \right]. \quad (3.5)$$

Although the imposed boundary condition on the flat plate is $v^- = 0$, the edge of the boundary layer also exists on the plate at $x = 0$. Hence, the vertical velocity at the edge, $v^+(x, \delta(x)) = c_v U_0/\sqrt{Re_x}$, also exists at $x = 0$. This would indicate that $v^+(0, 0)$ becomes infinite. However, the limit of $v^+(0, 0)$ should be a finite, non-zero value and with a net flux occurring through the leading edge point. Recognizing that the horizontal velocity $U_0$ is also incident to the leading edge we arrive at the following form of a mass balance at this point:

$$v_s(0) = U_0 - v^+(0, 0). \quad (3.6)$$

If we take $v^+(0, 0) = c_v U_0$, then we obtain $v_s(0) = U_0(1 - c_v)$ as a boundary condition. This means the horizontal velocity $U_0$ is partly ‘deflected’ or displaced into the vertical velocity $c_v U_0$, while the remainder is entrained into the leading edge point.

Figure 3(b) shows a set of computed wall shear stress $\tau_w$ from the vortex-entrainment sheet model and compares favorably with the Blasius result: $\tau_{w,B}/(\rho U_0^2) = 0.33/\sqrt{Re_x}$. The stress is still singular at $x = 0$ since the tangential velocity discontinuity of $u^- - u^+ = 0 - U_0$ remains. The inset plot shows that the relative error between the two calculations is bounded below 2% for $Re_x > 100$, which is consistent with our assumed region of validity [$Re_x \approx 100 \gg [(c_s c_m)^2 \approx 6]$.

4. The outer fluid solution

This section considers the mathematical formulation for the problem of determining the outer fluid velocity $u$. We discuss two different formulations, namely a Laplace equation and a boundary integral formulation. For a given problem, there may be certain advantages to using one or the other or a combination. For simplicity, we will now assume a two-dimensional flow. Hence, the vortex-entrainment sheet becomes a one-dimensional surface with $\mathbf{b}$ becoming the constant out-of-plane vector $\mathbf{k}$ so that the fluid and sheet velocities have zero $\mathbf{b}$ component and those relative to the $\mathbf{s-n}$ basis are $u = (u_s, u_n)$ and $v = (v_s, v_n)$. Likewise $v = \gamma \mathbf{k}$ and we can deal with the vortex sheet strength as a scalar quantity $\gamma(s,t) = u_s^- - u_s^+$. The section concludes with a suggested solution algorithm of the dynamically coupled system.

4.1. The Laplace equation

The problem is defined by specifying the governing equation and boundary conditions. The conjugacy of the harmonic potential $\phi$ and the stream function $\psi$ allows two equivalent formulations of the problem since both functions satisfy a Laplace equation. Here, we move forward with the $\phi$ framework, but provide a summary showing the duality with the $\psi$ framework. Some useful definitions relating $\Gamma$ and $Q$ to $\gamma$ and $q$ as well as to $\phi$ and $\psi$ may be found in Appendix B.
Let the open fluid domain be \( D \) with boundary/closure defined as the union \( \partial D = S_i \cup S_v \cup C_\infty \) (figure 4(a)), where \( S_i \) and \( S_v \) correspond to the vortex-entrainment sheets representing the surface with a sharp edge and the free sheet shed from the edge, respectively, and \( C_\infty \) is a large arbitrary boundary ‘at infinity’ that does not represent a surface of discontinuity. Usually, the boundary condition on \( S_i \) and \( S_v \) is prescribed as continuity of normal velocity, namely \((\partial \phi / \partial n)^+ = (\partial \phi / \partial n)^-\). Under this condition no entrainment occurs, and for a solid surface the normal velocity of that surface, \( U_n = (\partial \phi / \partial n)^\pm \), is specified.

However, the vortex-entrainment sheet is characterized by jumps in both the tangential and normal velocity components across the sheet. Hence, the values of \((\partial \phi / \partial n)^\pm\) and thus the entrainment rate \( q \) are assumed as given boundary conditions (however, see end of §4.3) . The special conditions for the sharp edge will be discussed in §5. With a fluid of infinite extent, the specification of the flow there serves as a boundary condition on \( C_\infty \), which may be expressed using an asymptotic form of the potential (Batchelor 1967, chap. 2.10). The statement of the problem thus far is:

\[
\nabla^2 \phi = 0 \quad x \in D \tag{4.1}
\]

\[
\phi(x) - \phi_\infty(x) \rightarrow C_\phi + \frac{\Gamma_\infty \theta}{2\pi} - \frac{Q_\infty}{2\pi} \log(r) + O(r^{-1}) \quad x \in C_\infty \tag{4.2}
\]

\[
\left. \frac{\partial \phi}{\partial n} \right|^\pm = u_n^\pm(x) \quad x \in S_i \cup S_v \tag{4.3}
\]

where \( C_\phi \) is a constant, \( \phi_\infty \) is an external potential representing a non-zero flow at infinity, \( \Gamma_\infty \) and \( Q_\infty \) are, respectively, the net circulation along \( C_\infty \) and the net flux across \( C_\infty \) into \( D \). In the case of a semi-infinite geometry \( S_i \) intersects \( C_\infty \) at infinity where the potential asymptotically becomes compatible with (4.2).

Now consider a contour \( C \) residing entirely within the analytic fluid region and bounding the open domain \( D \) as defined in figure 4(b). By mass conservation the total flux \( Q_C \) across \( C \) must be zero. This yields a ‘compatibility condition’ required for a solution \( \phi \) to this Neumann problem to exist (e.g. Stakgold 1968). Similarly, by circulation conservation (Kelvin’s theorem) the total circulation \( \Gamma_C \) around \( C \) must also be zero.
that need not decay at infinity. When all sheets:

\[ \text{Boundary condition:} \quad Q_C = -\oint_C \frac{\partial \phi}{\partial n} ds = 0 \]
\[ \quad \Gamma_C = \oint_C \mathbf{u} \cdot \hat{n} ds = \oint_C \frac{\partial \phi}{\partial s} ds = 0 \quad (4.4) \]

where \( s-n \) is an orthogonal coordinate system along \( C \). The \( Q_C \) equation can be derived from the Laplace equation itself upon using the divergence theorem. In general, the integration path is allowed to cross itself at the irregular point \( x_o \) in order to trace both sides of \( S_i \). When \( S_i \) represents a solid, no-slip surface this is not necessary because \( \nabla \phi = \mathbf{U} \) on the inside of \( S_i \). Since \( \partial \phi / \partial n \) is given everywhere on the boundary, then \( \mathbf{u}(x) = \nabla \phi \) is uniquely determined throughout the domain \( D \).

Lastly, in the \( \psi \) problem formulation the normal derivative \( \partial \psi / \partial n \) is given everywhere on the boundary and corresponds to the tangential velocity component. The conditions in (4.4) can be written with derivatives of \( \psi \) via the Cauchy-Riemann relations. Table 1 summarizes the two different frameworks, which are essentially two different, yet unique decompositions (Bhatia et al. 2013).

### 4.2. Boundary integral formulation

Now we write the outer flow \( \mathbf{u} \) as a complex conjugate velocity field \( \mathbf{w} = \mathbf{u} - iv \) and the contribution from the vortex-entrainment sheet is given by a Cauchy-type boundary integral. Jones (2003) obtained an elegant solution in this manner for the case of a moving flat plate with no entrainment by representing both the free and plate-bound vorticity as conventional vortex sheets. Namely, the two vortex sheets on each side of the plate were combined into a single sheet by coupling the plate and fluid velocities via a tangential boundary condition representing no-slip.

First, the velocity induced by the vortex-entrainment sheet in the two-dimensional version of (2.15) is expressed in a convenient complex form. Namely, \( \mathbf{x} \) is replaced with \( z = x + iy \) and the sheet position \( \mathbf{x}_s \) with \( \zeta = x_s(s) + iy_s(s) \) where \( s \) is the arclength coordinate. Introducing \( \theta(\zeta) \) as the angle measured from the horizontal to the local unit tangent vector on the sheet we have \( \partial \zeta / \partial s = e^{i\theta(s)} \). The conjugate velocity \( w(z) = u - iv \) at an analytic fluid point can then be written as:

\[ w(z) - w_\infty(z) = \frac{1}{2\pi i} \int_S \frac{\gamma(s) - i q(s)}{z - \zeta(s)} dz = \frac{1}{2\pi i} \int_S \frac{\chi(\zeta)}{z - \zeta} d\zeta, \quad (4.5) \]

where \( \chi(\zeta) = (\gamma - i q)e^{-i\theta} \) is the complex sheet strength and the integration is over all sheets: \( S = S_i \cup S_v \). Also, \( w_\infty(z) = d\phi_\infty / dz \) is the velocity of an external flow that need not decay at infinity. When \( z \to \zeta \) this singular integral is a generalized
Birkhoff-Rott equation for the vortex-entrainment sheet with complex strength. Following Muskhelishvili (1946) the left and right sides of the sheet are relative to an observer traversing the sheet in the positive direction of integration in (4.5) and the limits of any quantity approached from the left and right will be denoted with (+) and (−) superscripts, respectively. For any point on a sheet besides the sharp edge at $\zeta_o$ we have the following by the smooth-arc Plemelj formulæ (see comment in Appendix C):

$$\begin{align*}
\chi(\zeta) &= (w^- - w^+) = (\gamma - iq)e^{-i\theta} \\
w(\zeta) &= \frac{1}{2}(w^+ + w^-) = (\overline{u_n} - i\overline{u_{\infty}})e^{-i\theta}
\end{align*}$$

for $\zeta \in S_i \cup S_v \setminus \zeta_o$, \hspace{1cm} (4.6)

where $\overline{u}_n$ and $\overline{u}_{\infty}$ are the averages of the fluid velocity components. With this formulation it is easy to see how the sheet strength $\chi(\zeta)$ relates to the jump in the complex potential $\Phi = \phi + i\psi$ across the sheet (see Appendix B):

$$\begin{align*}
\Phi^- - \Phi^+ &= (\phi^- - \phi^+) + i(\psi^- - \psi^+) = \Gamma - iQ \\
\frac{\partial}{\partial s}(\Phi^- - \Phi^+) &= (u^-_n - u^+_n) - i(u^-_n - u^+_n) = \gamma - iq.
\end{align*}$$

Note that $\partial[\Phi]/\partial n = i(\partial[\Phi]/\partial s)$, again highlighting the conjugacy of the problem.

Now, the normal boundary condition (4.3) on a sheet $S$ can be expressed by:

$$\text{Re}\left\{\frac{1}{2}(w^+ + w^-)ie^{i\theta}\right\} = \frac{1}{2}(u^+_n + u^-_n) = \overline{u}_n$$

(4.9)

where each of the functions is evaluated at a position $\zeta$ on the sheet. Since we have supposed the normal velocities $u_{n}^\pm$ to be given, and thus the entrainment strength $q$ to be known, the boundary condition can be rearranged as:

$$f(\zeta) \equiv \text{Re}\left\{\frac{e^{i\theta(\zeta)}}{2\pi} \int_{S} \frac{\gamma(\xi)e^{-i\theta(\xi)}d\xi}{\zeta - \xi}\right\} = \overline{u}_n(\zeta) - w_{\infty,n}(\zeta)$$

$$+ \text{Re}\left\{\frac{e^{i\theta(\zeta)}}{2\pi} \int_{S} \frac{iq(\xi)e^{-i\theta(\xi)}d\xi}{\zeta - \xi}\right\}$$

(4.10)

where $\xi \in \mathbb{C}$ is a dummy integration variable along the sheet. This equation is to be solved for $\gamma(\zeta) = \gamma(s)$. The inversion formula and subsequent manipulations used by Jones (2003) are specific to the case of a plane boundary with $\theta(\zeta)$ constant along the sheet. As such, (4.10) is more amenable to numerical solution via, for example, expansion of $f(\zeta)$ as a Chebyshev series, which was also done by Jones as well as others since the series converges rapidly for smooth functions (Alben 2008).

4.3. Solution algorithm

We now present a brief solution algorithm for the coupled system of the vortex-entrainment sheet and the outer fluid. First, the outer flow $u$ is solved according to the specified normal boundary conditions with either the Laplace formulation (4.1)-(4.3) or the boundary integral formulation (4.5)-(4.10). For a free sheet we assume that it acquires circulation only as a result of having been shed from a body. Hence, $\gamma$ and $q$ will be known for these sheets. For a surface-bound sheet, the solution of the Laplace formulation will yield the velocity on the fluid side of the sheet, say $u^+$. Then, the additional boundary conditions on the surface side of the sheet are imposed to yield the sheet strengths. More specifically, we have $u^- = U$ with $U$ as the known surface velocity and thus $u^- - u^+ = \gamma \hat{s} + q\hat{n}$. The imposed conditions will typically be the no-slip and no-penetration conditions for a solid body, but in general we may model actual fluid slip on or porosity of the surface.
When the boundary integral formulation is used the specification of \( u_n^- = U_n \) on the surface side of the sheet will uniquely solve the flow inside that surface. Then the solution for \( \gamma \) will correspond to the \( u_s^- - u_s^+ \) of these two flows. To obtain the desired vortex sheet strength, the tangential boundary condition can be imposed by using the relative velocity \( (\overrightarrow{u}_s - U_s) \) in the evolution equation (2.24) as was done by Jones (2003). For a solid body, the harmonic potential \( \phi_b \) inside the body will be such that \( \nabla \phi_b = \mathbf{U} \); this is most easily seen in the reference frame of the body or for a stationary body where \( \phi_b = \text{const.} \) and thus the flow inside vanishes (Lamb 1945).

Equipped with \( \mathbf{u}, \gamma \) and \( q \) the solution for the sheet flow quantities can be obtained. The sheet mass and momentum equations (2.6)-(2.7) along with the evolution equation (2.24) represent 4 equations for the 4 unknowns \( \rho_s, \mathbf{v} = (v_s, v_n) \) and \([p]\); in three dimensions the system is augmented by the third component \( v_b \) and the corresponding vortex sheet strength component \( \gamma_s \). Again, these equations are coupled through the pressure jump and the entrainment strength. An initial condition at \( t_0 \) is required for the intrinsic flow quantities \( \gamma_s, \mathbf{v} = w + v_n \hat{n} \). In most cases, such as a flow started from rest, these initial conditions will simply be \( \gamma_s(x_s, t_0) = w(x_s, t_0) = 0 \) since no entrainment will have occurred yet anywhere in the domain.

Lastly, we note that it is theoretically possible to obtain both the vortex and entrainment sheet strengths if supplemental conditions are given elsewhere in the domain. In particular, if we are given the full velocity vector on \( C_\infty \) ‘at infinity,’ recall figure 4(a), then we may remove the requirement of \textit{a priori} specification of the normal boundary conditions on the sheets. By (4.2) and (4.5) we obtain:

\[
\begin{align*}
    w(z) - w_\infty(z) &= \left( \frac{\Gamma_\infty - iQ_\infty}{2\pi i R} \right) e^{-i\theta(z)} = \frac{1}{2\pi i} \int_S \frac{\gamma(s) - iq(s)}{z - \zeta(s)} ds \\
    &= \frac{1}{2\pi i} \int_S [\gamma(s) - iq(s)] ds, 
\end{align*}
\]

for \( |z| = R \to \infty \). The above integral is non-singular since \( z \notin S \) and we note that by writing \( z - \zeta(s) \approx Re^{i\theta(z)} \) then:

\[
\begin{align*}
    \Gamma_\infty - iQ_\infty &= \int_S [\gamma(s) - iq(s)] ds, 
\end{align*}
\]

which states that the circulation and flux in the sheets is equal to that at infinity or that the total of these quantities in the domain \( D \) is zero; recall (4.4). This condition makes the inversion formula of the singular integral equation in (4.10) unique (Jones 2003). The boundary conditions at infinity, for example by specification of \( \Gamma_\infty \) and \( Q_\infty \) on \( C_\infty \) would then determine the sheet strengths. However, the determination of \( \gamma - iq \) from (4.11) is the solution to a class of inverse problems, which are usually ill-conditioned. For this reason we do not pursue this option further.
5. Separation at the sharp edge

We now shift our focus to applying the vortex-entrainment sheet to the problem of separation at a sharp edge \( x_o = \zeta_o \). While the normal vector \( \hat{n} \) is uniquely defined on the wedge faces, \( x \neq x_o \), it is multivalued at the sharp edge. In order to have a well-posed problem this ambiguity must be removed. To see this, consider a wedge of interior angle \( \beta \pi \) with a rounded off corner of radius \( \varepsilon \ll 1 \) and with an irregular point existing somewhere in the domain as in figure 5. When the irregularity remains inside the surface (case i), \( \hat{n} \) varies continuously between the two limits \( \hat{n}^\pm \). When the irregularity is on the boundary of the ‘dulled’ edge (case ii), this point must be omitted to maintain an analytic fluid domain. As such, \( \hat{n} \) varies through values with a total change in argument of \( \pi \) and may have a direction outside those of \( \hat{n}^\pm \). As \( \varepsilon \to 0 \) the irregularity and geometric singularity merge (case iii), and not only does \( \hat{n} \) become ambiguous, but the change in argument is now \( 2\pi - \beta \pi \). Hence, the boundary normal vector at the sharp edge, or equivalently the shedding angle, must be defined such that (4.3) is compatible with the interaction of the flows approaching the corner. This process is governed by the intrinsic flow inside the sheets as is shown next. Unlike the case of a cusped edge the shedding angle for the wedge geometry will not, in general, be tangential to either face.

The sheet \( S_i \) is divided in two: \( S_1 \) and \( S_2 \) representing the wedge faces as depicted in figure 6(a). The subscripts (1), (2) and (v) will refer to quantities associated with \( S_1 \), \( S_2 \) and \( S_v \), respectively. The complex sheet strengths are then \( \chi_1 \), \( \chi_2 \) and \( \chi_v \), and the arguments of these sheets are \( \theta_1 \), \( \theta_2 \) and \( \theta_v \). The entrainment strength is known everywhere except at the sharp edge because the normal direction there is as yet undetermined. Imposing the normal boundary condition on the wedge faces will yield the corresponding vortex sheet strengths \( \gamma_1 \) and \( \gamma_2 \). It is assumed that the source of the vorticity in the free sheet \( S_v \) results solely from the merging/shedding of the two bound sheets \( S_1 \) and \( S_2 \). Hence, the remaining unknowns to be determined are \( \gamma_v \), \( q_v \) and \( \theta_v \).

First, we address the potential outer flow quantities relating to \( u \) and then discuss the flow \( v \) inside the sheets. By definition the sheets are characterized by the limiting values just outside the sheets. Hence, we have constraints to impose to ensure the consistency of \( S_1 \) and \( S_2 \) merging into \( S_v \). Figure 6(b) shows the tangential and normal components of the left and right limit velocities for each sheet. Next, we define \( w_v^\pm = (u_{s,v}^\pm - i u_{n,v}^\pm) e^{-i \theta_v} \) as the velocities on either side of \( S_v \). While \( w \) (i.e. \( u \)) may be discontinuous across a sheet it is still required to be smooth when \( z^\pm \in D \to (\zeta \in \partial D) \) since \( w \) is sectionally holomorphic (Muskhetlishvili 1946). This means that the \( w_v^\pm \) are related to the velocities.
just outside of $S_1$ and $S_2$ as:

$$w^+_v e^{i\theta_v} = (u^+_{s,1} - i u^+_{n,1})e^{-i(\theta_1 - \theta_v)}, \quad w^-_v e^{i\theta_v} = (u^-_{s,2} - i u^-_{n,2})e^{-i(\theta_2 - \theta_v)}. \quad (5.1)$$

The no-slip and no-penetration boundary conditions on the solid wedge faces can now be imposed since $u^-_{s,1}$, $u^+_{s,2}$, $u^-_{n,1}$ and $u^+_{n,2}$ will be known from the velocity of the wedge. Adding these components to the above equations, then subtracting and taking real and imaginary parts yields:

$$\gamma_v = u^-_{s,v} - u^+_{s,v} = \gamma_1 \cos(\theta_1 - \theta_v) + \gamma_2 \cos(\theta_2 - \theta_v)$$

$$q_v = u^-_{n,v} - u^+_{n,v} = \gamma_1 \sin(\theta_1 - \theta_v) + \gamma_2 \sin(\theta_2 - \theta_v) + [q_1 \cos(\theta_2 - \theta_v) + q_2 \cos(\theta_2 - \theta_v)]. \quad (5.2)$$

That $\gamma_v$ depends on $(q_1, q_2)$ and $q_v$ on $(\gamma_1, \gamma_2)$ is expected from the coupled sheet strength evolution equations in (2.24). Using the definition of the complex sheet strength in (4.6) these equations can be repackaged as:

$$\chi_v = \chi_1 + \chi_2. \quad (5.4)$$

In Appendix C it is shown that this requirement is precisely the condition that removes the singularities in the velocity induced at the sharp edge. Hence, (5.4) might as well be called the Kutta condition. However, note that the sheet strengths are not arbitrary, but are related to the balance of normal momentum and the pressure jump as discussed in §2.3 and referred to as the ‘neutralization of the singular pressure gradient’ in §2.4.

Now consider the intrinsic flows $v$ inside the sheets. Here we use vector notation to reiterate that $v$ is not part of the outer flow, but is confined to the sheet embedded within this flow. When the tangential flows $v_{s,1}$ and $v_{s,2}$ in the bound sheets merge at the edge point there must be no normal impulse relative to the shedding free sheet $S_v$. For otherwise there would be flow out of the sheet. The impulse represents the instantaneous or ‘ballistic’ merging of $S_1$ and $S_2$. This dynamic condition is expressed as:

$$0 = [(\rho_{s,1}v_{s,1}) \hat{s}_1 + (\rho_{s,2}v_{s,2}) \hat{s}_2] \cdot \hat{n}_v. \quad (5.5)$$

Noting from figure 6(a) that $\theta_2 = \beta \pi/2$ and $\theta_1 = -\beta \pi/2$, then (5.5) yields an expression for the shedding angle as:

$$\tan \theta_v = A \tan(\beta \pi/2), \quad A = \left( \frac{\rho_{s,2}v_{s,2} - \rho_{s,1}v_{s,1}}{\rho_{s,2}v_{s,2} + \rho_{s,1}v_{s,1}} \right). \quad (5.6)$$

Next, we have the mass and momentum boundary conditions of the merging process, which state that the $\rho_s$ and $\rho_s v_s$ from each of the bound sheets are carried into the free sheet by their respective convective velocities $v_s$:

$$\rho_{s,v}v_{s,v} = \rho_{s,1}v_{s,1} + \rho_{s,2}v_{s,2}, \quad \rho_{s,v}v_{s,v}^2 = \left( \left(\rho_{s,1}v_{s,1}^2 \right) \hat{s}_1 + (\rho_{s,2}v_{s,2}^2) \hat{s}_2 \right) \cdot \hat{s}_v. \quad (5.7)$$

With the intrinsic flow quantities $\rho_{s,j}$ and $v_j$ ($j = 1, 2$) known from the solution of the coupled system, then (5.2), (5.3) and (5.6) represent three equations for the three unknowns $\gamma_v$, $q_v$ and $\theta_v$. Then (5.7) gives boundary conditions to calculate the flow into $S_v$ as it is shed.

5.1. Discussion on special cases of entrainment

In this subsection we discuss some sets of simplified circumstances pertaining to separation at a sharp edge that are of practical interest. When $\beta = 0$ then (5.6) simply
gives $\theta_v = 0$ irrespective of the intrinsic flow $\mathbf{v}$ and thus $[p]$. Of course this corresponds to the usual tangential shedding result for a cusped edge.

Now consider zero entrainment everywhere so that there is no intrinsic flow $\mathbf{v} = 0$ or mass $\rho_s = 0$ in the sheets, namely the conventional vortex sheet. Equation (5.6) then becomes indeterminate and a new condition for the shedding angle is required. This is obtained from (5.3), which with each $q_j = 0$ reduces to:

$$\gamma_1 \sin(\beta \pi / 2 + \theta_v) = \gamma_2 \sin(\beta \pi / 2 - \theta_v).$$

(5.8)

There is the trivial solution with $\gamma_1 = \gamma_2 = 0$ meaning no flow at the edge. The only other solutions that put $\theta_v$ in a physical range are $\theta_v = \pm \beta \pi / 2$ and with the respective consequences of $\gamma_1 = 0$ or $\gamma_2 = 0$. This corresponds to a tangential shedding angle and that the flow is stagnated on one side of the vortex sheet depending on the sign of shed circulation. Hence, as mentioned in §2.4 the vortex sheet from one wedge face convects off the surface with no contribution from the opposing face. These features are consistent with the results of Pullin (1978) in which there was zero entrainment.

Lastly, consider the onset of motion at $t = 0$ when no vorticity has been shed from the sharp edge and $S_v$ does not yet exist. As used by Rott (1956) and Pullin (1978) the attached potential flow responsible for separation is obtained as the leading order term in an expansion of the complex potential near the sharp edge: $\Phi \sim (z - \zeta_0)^n$ where $n = 1/(2 - \beta)$ for the wedge of angle $\beta \pi$. The imposed pressure gradient at infinity that drives this flow suggests the windward wedge face will ‘feel’ the pressure at $t = 0^+$ before the leeward face. Hence, we would expect more significant entrainment to occur on the windward face, say $S_2$, so that $v_{2,s} \gg v_{1,s}$ and (5.6) gives $A \approx 1$ and $\theta_v \approx \beta \pi / 2$. This is again consistent with the small-time similarity solution. In a real flow the ensuing roll-up would be sure to quickly induce a pressure and so entrainment on the leeward face thus changing $\theta_v$.

6. Example calculations

This section presents two example calculations to demonstrate the vortex-entrainment sheet concept as applied to separation at a sharp edge. For simplicity we make some further assumptions on the dynamics for a free sheet. This is accomplished by enforcing a relation between $\mathbf{v}$ and $\mathbf{u}$. To this end, we recognize that the conventional vortex sheet has the Birkhoff-Rott equation as its evolution equation, which says the sheet moves with the average fluid velocity. Motivated by this, we simply set $\mathbf{v} = \overline{\mathbf{u}}$ and we then have the generalized Birkhoff-Rott equation (4.5) to convect the free sheet.

To see the dynamical consequences of our assumption we apply them to the mass and momentum equations (2.6)-(2.7) to give:

$$\frac{\partial \rho_s}{\partial t} + \frac{\partial}{\partial s}(\rho_s v_s) = \rho q, \quad \rho_s \frac{D_s v_s}{D_t} = [\tau], \quad \rho_s \frac{D_s v_n}{D_t} = -[p].$$

(6.1)

We see that the normal momentum equation remains unchanged regardless of the value of the entrainment strength $q$. Hence, the pressure jump is exactly that required to move the mass in the sheet with $v_n = \overline{u_n}$. For a free sheet we have $\tau^+ = \tau^- = 0$ since the outer flow is potential and thus $[\tau] = 0$ means the tangential motion of particles in the sheet move materially.

6.1. Setup of computed simulations

First we present the problem setup for a case with zero entrainment that is given in §6.2. A case with non-zero entrainment is then considered in §6.3 and the required
modifications to the problem setup are also discussed there. The geometry of each problem is a wedge of infinite extent.

We use a combination of the Laplace and boundary integral formulations. The wedge is stationary and a boundary condition is imposed at infinity to drive the flow in the form of a known harmonic potential function \( \phi_\infty(r, \theta, t) = \text{Re} \{ \Phi_\infty(z, t) \} \) that represents the attached flow. Since there is no entrainment \( S_1 \) and \( S_2 \) have no-penetration conditions and \( S_v \) has a continuity of normal velocity condition. Moreover, by construction the potential \( \phi_\infty \) satisfies homogeneous Neumann boundary conditions on the wedge faces. Hence, after each new free vortex sheet segment is shed, we need only to account for the normal velocity induced by \( S_v \) on the walls. We can then arrange the problem for a potential \( \phi_w \), which represents the image system of the free vortex sheet inside the wedge. This \( \phi_w \) satisfies:

\[
\nabla^2 \phi_w = 0 \quad \text{for} \quad r > 0, \quad \theta \in [-\theta_w, \theta_w] \tag{6.2}
\]

\[
\frac{1}{r} \frac{\partial \phi_w}{\partial \theta} = -\text{Re} \{ i e^{i \theta} w_v(r, \theta) \} \quad \text{for} \quad r > 0, \quad \theta = \pm \theta_w \tag{6.3}
\]

\[
\phi_w \rightarrow 0 \quad \text{for} \quad r \rightarrow \infty, \quad \theta \in [-\theta_w, \theta_w] \tag{6.4}
\]

where \( \theta_w = \pi(1 - \beta/2) \) and \((\pm)\) corresponds to the angular coordinates of \( S_1 \) and \( S_2 \), and so \( w_v(r, \pm \theta_w) \) is the velocity induced by \( S_v \) on the wedge faces. For brevity we will write \( \zeta_w = re^{i\theta_w} \) for the coordinate on either \( S_1 \) or \( S_2 \).

In §5.1 we saw that for zero entrainment the shedding angle equation (5.6) becomes indeterminate and \( \theta_v \) is instead obtained by (5.3) with \( q_v = q_1 = q_2 = 0 \). Specifically \( \theta_v \) is necessarily tangent to one wedge face and the flow is stagnated at the apex on the other face. This second condition could be imposed by \( \partial \phi / \partial r = 0 \) on the \( S_1 \) side where \( \phi \) is the total potential. However, it is automatically satisfied by imposing \( \theta_v \equiv \beta \pi/2 \). We note that this is a result of having to impose the no through-flow condition on both \( S_1 \) and \( S_2 \) as well as the continuity of normal velocity on each side of \( S_v \), all of which exist at the same location. As a consequence of all these constraints, the strength of the new vortex sheet segment of \( S_v \) is \( \gamma_v = \gamma_2 \) by (5.2), which represents \( S_2 \) convecting off to form \( S_v \).

The calculation procedure is as follows. Given at time \( t_0 = 0 \) is the vortex sheet strength \( \gamma_w \) on \( S_1 \) and \( S_2 \) from \( \phi_\infty(\zeta_w, t_0) \) and an initialized position \( \zeta_v \) and vortex sheet strength \( \gamma_v \) of the free sheet \( S_v \). The flow is initialized by representing the first shed segment of \( S_v \) as a point vortex as was similarly done by Jones (2003). Then, for each \( k \geq 1 \) time step:

(i) Compute new velocity on wedge faces due to \( S_v \) as

\[
w_v(\zeta_v) = \frac{1}{2\pi i} \int_{\zeta_v - \zeta_v(s)} \frac{\gamma_v(s)ds}{\zeta_v - \zeta_v(s)} \tag{6.5}
\]

(ii) Solve the Laplace equation for \( \phi_w \) in (6.2) with boundary conditions (6.3)-(6.4).

(iii) Compute new \( \gamma_w \) along \( S_1 \) and \( S_2 \), i.e. \( \gamma_1(r) \) and \( \gamma_2(r) \), as:

\[
\gamma_w(\zeta_w) = \frac{\partial}{\partial r}(\phi_w(\zeta_w) + \phi_\infty(\zeta_w)) + \text{Re} \{ w_v(\zeta_w)e^{i\theta(\zeta_w)} \} \tag{6.6}
\]

(iv) Set \( \gamma_v = \gamma_2(r = 0) \) and \( \theta_v = \beta \pi/2 \) of new sheet segment of \( S_v \).

(v) Compute new total induced velocity on free sheet \( S_v \) as

\[
w(\zeta_w) = \frac{\partial \zeta_v}{\partial t} = \frac{d\Phi_\infty}{dz} + \frac{1}{2\pi i} \int_{S_v} \frac{\gamma_v(s)ds}{\zeta_v - \zeta_v(s)} + \frac{1}{2\pi i} \int_{S_1 \cup S_2} \frac{\gamma_w(s)ds}{\zeta_v - \zeta_w(s)} \tag{6.7}
\]

(vi) Time integrate \( \partial \zeta_v / \partial t \) to advect free sheet \( S_v \) to \( t_{k+1} \).
Figure 7. (a) Comparison of streamlines (solid) and vortex sheet position (dashed) for the case with $m = 0$ and $\beta \pi = \pi/2$ ($= 90^\circ$) at $t = 1$ scaled to the similarity space $\omega = \xi + i\eta$. (b) The total and rate of shed circulation as a function of time. The inset shows the radial velocity along the upper wall, $u_r$, as a function of distance from the apex normalized by the separation bubble length $L_b$. The flow is stagnated at the apex and the reattachment point $r = L_b$ with reversed flow in the separation bubble.

In steps (v) and (vi) $\overline{\zeta}$, is the complex conjugate of the sheet position. Due to the simple geometry we use the method of images that maps the domain outside the wedge to a semi-infinite plane as $z^* = z^n$. This combines steps (i)-(iii) to obtain $\gamma_w$ directly. The integrals may be computed with a discrete sheet method or as a system of point vortices. Lastly, the free sheet $S_v$ is advected with a fourth-order Runge Kutta scheme.

6.2. Starting flow past an infinite wedge

As validation we first reproduce a result from the similarity solutions of Pullin (1978) for a starting flow over an infinite wedge of interior angle $\beta \pi$. There is zero entrainment everywhere for all time and the attached flow complex potential is $\Phi_\infty(z,t) = -iat^mz^n$ where $n = 1/(2 - \beta)$. We simulate a time-dependent flow beginning from $t = 0$.

Figure 7(a) shows a comparison of the streamlines and vortex sheet location for the case with $m = 0$ and $\beta \pi = \pi/2$. The computed simulation is at $t = 1$ and is scaled to the similarity space $\omega = \xi + i\eta$ (for more detail see Pullin 1978). There is good agreement between the visual character of the flow field. For a more quantitative comparison figure 7(b) plots the total and rate of shed circulation, again showing excellent agreement with the similarity laws given by Pullin. The inset shows the velocity along the upper wedge surface (i.e. $S_1$), which stagnates at the leeward side of the apex as expected from the boundary condition for the case of zero entrainment.

6.3. Oscillating flow with estimated entrainment

Technically the entrainment strengths are to be given as boundary conditions. However, we may estimate the entrainment $q_v$ of the free sheet by relaxing the condition of stagnated flow on the leeward side of the apex. In other words, a non-zero wall-tangential velocity is allowed just outside both $S_1$ and $S_2$ at the apex: $\gamma_1 \neq 0$ and $\gamma_2 \neq 0$. Each of these velocities then has a normal component relative to the free sheet and must merge to form $S_v$. Hence, the merging of the $S_1$ and $S_2$ sheets is a form of entrainment. This means a non-tangential shedding angle is possible and this is determined next.

These wall-tangent flows are in the outer potential flow $u$ just upstream of the sharp edge point and thus must be entrained through the sharp edge to become the intrinsic
flow \( v \). Therefore, since there cannot be flow through \( S_v \) the normal impulse relative to \( S_v \) of these incoming flows must be equal just as was the case for (5.5). Hence, we obtain an analogous equation for the shedding angle:

\[
\tan \theta_v = B \tan(\beta \pi/2) \quad B = \left( \frac{v_s^{-2} - v_s^+}{v_s^{-2} + v_s^+} \right). \tag{6.8}
\]

Although we have assumed \([\rho]\) = 0 we could accommodate a stratified flow by the substitution \( u^\pm \rightarrow \rho^\pm u^\pm \). The above equation is equivalent to the one used by Xia & Mohseni (2017); the two relations are related by the trigonometric identity \( \arctan(x) = \arccos(1/\sqrt{1+x^2}) \). However, there it was not recognized that this necessarily requires entrainment into the shed sheet.

Since \( \theta_v \) is now determined by (6.8) instead of prescribed as \( \beta \pi/2 \), then step (iv) of the procedure given in §6.2 must be updated as:

(iv)' Compute \( \gamma_v \) and \( q_v \) from (5.2) and (5.3) from \( \gamma_1 \), \( \gamma_2 \), \( \theta_v \) with \( q_1 = q_2 = 0 \).

In the same way that a segment of vortex sheet must be of finite length to possess circulation, \( \Gamma = \int \gamma ds \), so too must an entrainment sheet segment have a finite length to have an entrainment rate, \( Q = \int q ds \). By confining the entrainment over some small sheet segment \( \Delta s \) to the point at the sharp edge means we are essentially approximating a net entrainment rate \( Q \) as

\[
Q = \int_{\Delta s} q_v ds \approx q_v \Delta s. \tag{6.9}
\]

Then the amount of mass of fluid with density \( \rho \) that is entrained is \( \Delta m \approx \rho Q \Delta t \). Hence, the sheet mass density \( \rho_{s,v} = \Delta m/\Delta s \) put into \( S_v \) in a given amount of time is:

\[
\rho_{s,v} \approx \rho q_v \Delta t, \tag{6.10}
\]

which is equivalent to integrating the mass equation in (6.1) over \( \Delta t \). Therefore, with a constant \( \Delta t \) and \( \rho \) the behavior of \( \rho_{s,v} \) and \( q_v \) are synonymous with the understanding that \( q_v \) exists only at the edge and puts \( \rho_{s,v} \) in the sheet which then convects downstream. In other words, the time trace of \( q_v \) is analogous to the trace of \( \rho_{s,v} \) along the sheet arclength at a given time.

Having physically explained the estimated entrainment through the sharp edge, we now compute a solution that is similar to the wake created by the trailing edge of an oscillating airfoil. The driving potential flow is given by

\[
\Phi(z,t) = A_1(t)(z - \zeta_o)^n + A_2(t)(z - \zeta_o)^{2n}, \tag{6.11}
\]
Figure 9. Computed quantities for the simulated flow in figure 8. (Top) The entrainment sheet strength at the sharp edge, (middle) shedding angle $\theta_v$ where the dashed lines are the tangential limits $\beta \pi/2 = \pm 8^\circ$, and (bottom) vortex sheet strength at the edge. All are plotted against time. The solid circles and open triangles respectively mark where the entrainment strength is strongest and weakest.

where $A_1(t) = A_o \cos(2\pi ft)$ is the oscillating component and $A_2 = U_\infty$ represents the translational velocity of the airfoil. To provide a qualitative comparison we matched the oscillation frequency from one of the dye visualization experiments of Koochesfahani (1989). The experiments used a NACA 0012 airfoil with chord-based Reynolds number $Re_c = 11,400$ and $2^\circ$ oscillation amplitude and frequencies $f = 4, 5$ and $6$ Hz. We ran a simulation with $f = 6$ Hz and where $q_v$ was estimated as described above. Due to the more complicated sheet shape, we used a point vortex method to approximate the sheet.

Figure 8(a) plots a comparison of the simulated wake structure at $t = 2$ and shows reasonably good agreement with the experiment. Figure 8(b) shows a zoomed in view of the sharp edge and plots the free sheet shape at several times through a period of oscillation. We can clearly see the smooth variation of the shedding angle $\theta_v$. The computed entrainment sheet strength $q_v$, shedding angle $\theta_v$ and vortex sheet strength $\gamma_v$ are the top, middle and bottom plots in figure 9, respectively. First, we see that $q_v$ oscillates with twice the frequency as the oscillating flow, but is shifted upward by a positive mean, $q_v \approx 0.159$, such that $q_v(t) > 0$. This is to be expected since entrainment is a viscous process that cannot be reversed such that the fluid is returned to an irrotational state. The shedding angle $\theta_v$ shows the expected oscillatory behavior within the two tangential shedding limits of $\pm \beta \pi/2$ and about a zero mean, namely the wedge bisector and likewise for the vortex sheet strength.

The solid circles mark when $q_v$ is strongest and the open triangles mark when $q_v$ is weakest. Although we have oscillated the flow, we will now speak as if the trailing edge were oscillating instead. Hence, entrainment is strongest when the trailing edge is passing through the zero incidence position. At these times the edge is moving with its highest velocity and changes sign of acceleration. This deceleration is known to cause an impingement of the flow onto the leeward side of the wedge and a corresponding significant normal force (DeVoria & Ringuette 2013). This large pressure/normal force on the wedge is then responsible for the strong entrainment. The vortex sheet displays the
opposite behavior showing maxima and minima when the wedge changes sign of velocity at the largest amplitudes of the oscillation. This conjugacy between the vortex sheet and entrainment sheet, or in general tangential and normal quantities, is a feature that we have repeatedly observed throughout this work.

7. Concluding remarks

This study proposed the vortex-entraiment sheet as a model of viscous boundary and shear layers in three-dimensional flow. The sheet differs from the conventional vortex sheet by allowing mass and consequently momentum in the sheet. The process of entrainment is the mechanism that allows fluid to enter the sheet, thus endowing it with inertia. Hence, there is an intrinsic flow confined to the tangential manifold defining the sheet position. This internal flow is dynamically coupled to the flow outside the sheet and may also have physical properties and phenomena that are different from the bulk fluid and that exist only on the sheet.

The physical concept motivating the entrainment sheet definition is the preservation of the mass within a finite-thickness viscous layer in the limit as the layer collapses to zero thickness. In other words, the sheet has a mass per-unit-area of sheet, \( \rho_s \). The essential characteristics of the sheet are the vectorial vortex sheet strength representing tangential discontinuities in the velocity, and the scalar entrainment sheet strength corresponding to a discontinuity of the normal velocity. The latter strength is obviously responsible for the entrainment and acts as a source of \( \rho_s \) relative to the flow inside the sheet.

The velocity field induced by the vortex-entraiment sheet is given by a generalized Birkhoff-Rott equation where the sheet has a complex strength with the real part corresponding to the vortex sheet as usual and the imaginary part to the entrainment sheet. However, this generalized equation is not necessarily the evolution equation for the sheet since it contains mass, which also allows it to support a pressure jump. Another major physical consequence of the sheet mass is that no additional Kutta condition is required as the flow remains finite due to a balance of normal momentum with the pressure jump. For the specific case of separation at a sharp edge, the shedding angle is dictated by the normal impulse of the intrinsic flows in the sheets merging at the edge to form the free sheet. These features represent stark differences from the conventional vortex sheet, which remains dynamically indistinct from the surrounding fluid. A massless vortex sheet, \( \rho_s \equiv 0 \), is essentially decoupled from a governing momentum equation and thus is said to ‘move with the fluid.’ For the same reason, an explicit Kutta condition at irregular points is required to replace the degenerate momentum equation for the conventional vortex sheet.

Throughout this work a theme of conjugacy between the vortex sheet/tangential quantities and the entrainment sheet/normal quantities was observed. We interpret this to mean the entrainment sheet is a natural extension of the conventional vortex sheet as a model of viscous layers. The vortex-entraiment sheet concept was applied to a few example calculations and showed encouraging results about representing the viscous phenomenon of entrainment with an inviscid model.

This paper is the first part in a planned two-part series. The second part will extend the vortex-entraiment sheet model to the case of separation on a smooth surface where the separation point is able to move relative to the surface. In addition, the problem of predicting the location of the separation point is considered.
Acknowledgments

We wish to acknowledge the partial support of the AFOSR and ONR in this work. We are also grateful for discussions with Professors D. I. Pullin and A. Leonard pertaining to an earlier version of this manuscript.

Appendix A: Singular parts of the curl and divergence

This appendix represents the discontinuity in $u$ across the sheet as a Heaviside function to obtain the singular parts of the curl and divergence operators as given in (2.12). The following derivation essentially parallels that given by Howe (2007) for a vortex sheet. Near the sheet the discontinuous fluid velocity can be expressed as:

$$u = H(n^+)u^+ + H(n^-)u^-$$

(7.1)

where $n = 0$ specifies the sheet location and $H(n)$ is the Heaviside function with defining property $H(n > 0) = -H(n < 0) = H(n)$. Substituting this into the curl and divergence of the velocity gives:

$$\omega = \nabla \times u = \nabla H(n) \times (u^+ - u^-) \quad \Delta = \nabla \cdot u = \nabla H(n) \cdot (u^+ - u^-)$$

(7.2)

since $u^+$ and $u^-$ are incompressible, irrotational flows existing in the open region defined by the fluid domain $D$ whose closure contains the sheet. Substituting $\nabla H(n) = \delta(n)\hat{n}$ gives:

$$\omega = \hat{n} \times [u]\delta(n) = \gamma\delta(n) \quad \Delta = \hat{n} \cdot [u]\delta(n) = -q\delta(n).$$

(7.3)

where $\gamma(s, b)$ is a sheet-tangent vector giving the strength of the vortex sheet, and $q(s, b)$ is the strength of the entrainment sheet as given by (2.13) and (2.14).

Appendix B: Conjugate definitions

Let $s$ and $n$ be the tangential and normal sheet coordinates on a one-dimensional sheet immersed in a two-dimensional flow. The amount of circulation within and the flux into a sheet segment are:

$$\Gamma = \oint u \cdot dl = \int \gamma ds, \quad Q = -\oint u \cdot \hat{n} dl = \int q ds.$$

(7.1)

Using $u = \nabla \phi$ and the Cauchy-Riemann relations, the tangential and normal components of velocity, $u_s$ and $u_n$, are

$$u_s = \frac{\partial \phi}{\partial s} = \frac{\partial \psi}{\partial n}, \quad u_n = \frac{\partial \phi}{\partial n} = -\frac{\partial \psi}{\partial s}.$$

(7.2)

As such, the integrands in (7.1) can be written as exact differentials of $\phi$ and $\psi$, respectively. Therefore, integrating on each side of the sheet we have:

$$\Gamma = -(\phi^+ - \phi^-) = -[\phi], \quad Q = (\psi^+ - \psi^-) = [\psi].$$

(7.3)

Taking the partial derivative of these quantities along the tangential coordinate gives the vortex and entrainment sheet strengths, $\gamma$ and $q$, as:

$$\gamma = \frac{\partial \Gamma}{\partial s} = u_s^- - u_s^+, \quad q = \frac{\partial Q}{\partial s} = u_n^- - u_n^+.$$

(7.4)

It is worth reiterating that $Q > 0$, $q > 0$ corresponds to entrainment into the sheet as signified by the negative sign in the equation defining $Q$. 
Appendix C: Removal of singularities

Comment: In the theory of singular integral equations the sheet strength is conventionally defined as the jump in \( w \), that is \( \chi = [w] \) and the Cauchy kernel is \((\zeta - z)^{-1}\). However, in (4.5) and (4.6) we defined the sheet strength to be minus the jump so that the kernel is \((z - \zeta)^{-1}\). This was done so as to avoid introduction of a superfluous minus sign. Without loss of generality and to be consistent with the equations of Muskhelishvili (1946) and others, we here adopt the conventional definition.

This appendix examines the behavior of the complex velocity \( w(\zeta) \) as the sharp edge is approached, \( \zeta \to \zeta_o \). In general there are discontinuities in the strengths at \( \zeta_o \) since this point represents the confluence of the sheets \( S_1, S_2 \) and \( S_v \). As such, the sheet strength \( \chi(\zeta) \) does not satisfy the Hölder condition (Muskhelishvili 1946) exactly at \( \zeta_o \), and so the Plemelj formulae for a corner point do not apply. However, each sheet strength may be said to satisfy the Hölder condition on its respective closed arc \( S_1, S_2 \) and \( S_v \), in which case formulae are known for the behavior near the end points (see ibid).

The velocity induced by a given sheet at a point \( \zeta \) near the corner point is:

\[
\begin{align*}
    w(\zeta) &= \frac{1}{2\pi i} \int_{\zeta_o}^{\zeta} \frac{\chi(\xi)}{\xi - \zeta} d\xi = \frac{\chi(\zeta_o)}{2\pi i} \int_{\zeta_o}^{\zeta} \frac{d\xi}{\xi - \zeta} + \frac{1}{2\pi i} \int_{\zeta_o}^{\zeta} \frac{\chi(\xi) - \chi(\zeta_o)}{\xi - \zeta} d\xi \\
    &= \frac{\chi(\zeta_o)}{2\pi i} \log(\zeta - \zeta_o) + G(\zeta)
\end{align*}
\]

where \( G(\zeta) \) satisfies the Hölder condition near and at \( \zeta_o \). The velocities of this form induced by the individual sheets are physically meaningless as \( \zeta \to \zeta_o \), becoming logarithmically infinite, and the approaches from the left and right are not well defined. The singularity may be of a more general form if it is assumed that \( \chi(\zeta) = \chi^*(\zeta)/(\zeta - \zeta_o)^\alpha \) near the corner where \( \alpha \in \mathbb{C} \) is a complex constant. Although this may yield a finite \( w(\zeta_o) \), we take the physical significance of the sheet strengths \( \chi \) to imply that they too remain bounded, so that \( \alpha = 0 \) and the singularity is logarithmic.

Therefore, we must consider the collective behavior of the total induced velocity. Let \( \zeta_1, \zeta_2 \) and \( \zeta_o \) be positions on the sheets \( S_1, S_2 \) and \( S_v \) such that \( \epsilon = |\zeta_1 - \zeta_o| = |\zeta_2 - \zeta_o| = |\zeta_v - \zeta_o| \) are short equidistant lengths from the corner point \( \zeta_o \). Then:

\[
    w(\zeta_o) = \frac{\chi_1 + \chi_2 - \chi_v}{2\pi i} \log \epsilon + G(\zeta_o)
\]

where \( \chi_1, \chi_2 \) and \( \chi_v \) are each evaluated at \( \zeta_o \). Hence, the constraint \( \chi_v = \chi_1 + \chi_2 \) ensures that this logarithmic singularity in the velocity is removed, and is the same condition given in (5.4) of §5.

Now, the complex potential \( \Phi(z) \) is determined by the same Cauchy-type integral as the conjugate velocity \( w(z) \), so that \( \Phi \) has an analogous logarithmic singularity. This is the source of the inverse square-root singularities in the velocity as mentioned by Jones (2003). More specifically since \( w = d\Phi/dz \), then:

\[
\begin{align*}
    \frac{d}{d\zeta} \log |\zeta - \zeta_o| &\sim \frac{1}{|\zeta - \zeta_o|} = \frac{1}{\sqrt{(x - x_o)^2 + (y - y_o)^2}}. \\
\end{align*}
\]

where recall that \( \epsilon = |\zeta - \zeta_o| \). However, these singularities are also removed by the conditions imposed on the \( \chi_j \). This is due to the relationship between the doublet sheet strength and the vortex-entrainment sheet strength as \( \chi = \partial \lambda / \partial \zeta \). In fact, this is what allows the velocity to be written as:

\[
    w(z) = \frac{d\Phi}{dz} = \frac{1}{2\pi i} \int_S \frac{\chi(\xi)}{z - \xi} d\zeta = \frac{1}{2\pi i} \int_S \frac{(\partial \lambda / \partial s)ds}{z - \zeta(s)}.
\]
A. C. DeVoria and K. Mohseni

REFERENCES


