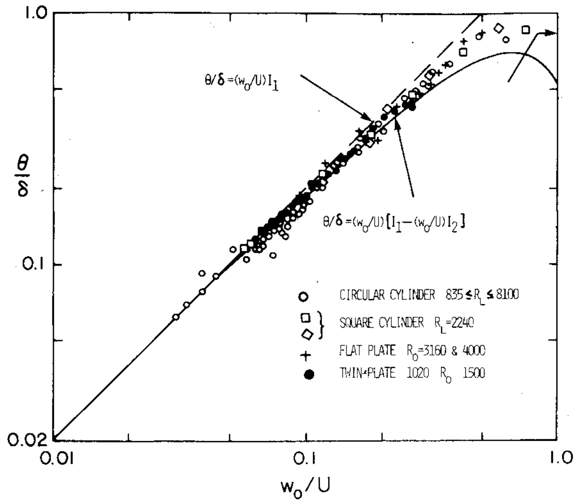

Fig. 3 A plot of θ/δ vs w_0/U. For $w_0/U\leq 0.05$, Eq. (4) and the linear approximation to it are indistinguishable.

Some relation can be expected to exist between the near-wake behavior and the boundary layers just before they leave the wake generators. The boundary layers near the trailing edge of both flat plates examined here were known to be turbulent; the boundary layers on the bluff bodies become unstable and turbulent soon after they separate. In view of this, it is quite surprising to find that θ/δ values for all wakes seem to approach, as $w_0/U\to 1$, the value ($\simeq 0.856$) appropriate to a flat-plate laminar boundary layer (with θ and δ both defined as done here for wakes); although the corresponding number for a turbulent boundary layer depends on the Reynolds number, it is at least an order of magnitude higher. As our measurements stopped short of extending all the way up to the wake generator only by a dozen or so momentum thicknesses, the last result should simply mean that, within a few momentum thicknesses after leaving the wake generator, the flow quickly readjusts as if the boundary layers leaving the wake generator were laminar, a curious behavior worth a closer examination!

At lower Reynolds numbers (R_θ of the order of 300 or lower), preliminary measurements behind flat plates showed an even more complex behavior in which the parameters W and Δ showed discontinuous jumps when plotted against w_0/U. Presumably, some of this is associated with the Karman vortex patterns that have an important influence on flow development at these Reynolds numbers.

Acknowledgments

All of the wake parameters reported here have been obtained by a reanalysis of the raw data collected by my colleagues and me at the Indian Institute of Science, Bangalore, over a period of about four years. My thanks go to Drs. A. Prabhu, P. K. Bhutiani, and Mr. Prasanna Kumar, whose measurements I have used, and to Prof. R. Narasimha, whose constant interest in the matter actually spurred the analysis.

References

[1]Townsend, A. A., *The Structure of Turbulent Shear Flow,* University Press, Cambridge, U.K., 1956.

[2]Prabhu, A., "Non-equilibrium Wakes," Ph.D. Thesis, Dept. of Aero. Eng., Indian Institute of Science, Bangalore, 1971.

[3]Bhutiani, P. K., "Effect of Free Stream Turbulence on the Development of a Turbulent Wake," M.E. Dissertation, Dept. of Aero. Eng., Indian Institute of Science, Bangalore, 1972.

[4]Prasanna Kumar, "A Study of Wakes Behind Two-Dimensional Bodies," M.E. Dissertation, Dept. of Aero. Eng., Indian Institute of Science, Bangalore, 1972.

[5]Chevray, R. and Kovasznay, L.S.G., "Turbulence Measurements in the Wake of a Thin Flat Plate," *AIAA Journal*, Vol. 7, Aug. 1969, pp. 1641-1643.

[6]Sreenivasan, K. R., "Data on Constant-Pressure Two-Dimensional Turbulent Wakes," Indian Institute of Science, Bangalore, Aero. Eng. Rept. 74FM5, 1974.

[7]Narasimha, R. and Prabhu, A., "Equilibrium and Relaxation in Turbulent Wakes," *Journal of Fluid Mechanics*, Vol. 54, July 1972, pp. 1-17.

[8]Sreenivasan, K. R. and Narasimha, R., "Equilibrium Parameters for Two-Dimensional Turbulent Wakes,".

[9]Prabhu, A. and Narasimha, R., "Turbulent Non-Equilibrium Wakes," *Journal of Fluid Mechanics*, Vol. 54, July 1972, pp. 19-38.

AIAA 81-4277

Transformation of the Equation Governing Disturbances of a Two-Dimensional Compressible Flow

E. J. Kerschen* and T. F. Balsa†
General Electric Corporate Research and Development Center, Schenectady, N.Y.

THIS Note concerns the propagation of small amplitude inviscid and unsteady disturbances on steady nonuniform mean flows. Problems of this type arise in a wide variety of fields, including aerodynamic noise, flutter, forced vibration, and buffeting of structures; for example, when a purely vortical disturbance is incident on an airfoil at *large* angles of attack. Indeed, the present discussion is restricted to the analysis of the distortion of vortical and entropic disturbances as they convect, with the mean flow, past a "bluff" object (flow separation is ignored) and to the generation of certain irrotational fields that permit the enforcement of boundary conditions on the body surface. A classical approach is to reconstruct the velocity field from the vorticity and then to obtain the pressure field from the former.[1] Goldstein[2] has developed a much simpler approach which requires the solution of a single inhomogeneous wave-like equation for the irrotational (roughly the acoustic) field. Because the mean flow is nonuniform, this equation has variable coefficients.

The purpose of this Note is to show that, when the behavior of the two-dimensional and compressible mean flow is approximated by the tangent gas relations,[3] the inhomogeneous wave equation can be transformed into a much simpler form involving only one variable coefficient. Further simplifications are possible when the mean flow is a small perturbation of a uniform stream.

We consider small amplitude disturbances superimposed on a steady, irrotational, compressible mean flow. Linearizing the equations of motion about the mean flow, and neglecting viscous effects, Goldstein[2] has shown that the perturbations are described by the following equations

$$u'=\nabla G+v \tag{1}$$

$$s'=b(X-iU_\infty t) \tag{2a}$$

Received April 6, 1981.

*Mechanical Engineer.

†Mechanical Engineer. Member AIAA.

$$v_i=\left(\frac{s'}{2c_p}\right)U_i+a(X-iU_\infty t)\bullet\frac{\partial X}{\partial x_i}\qquad i=1,2,3 \tag{2b}$$

$$p'=-\rho_0\frac{D_0G}{Dt} \tag{3a}$$

$$\rho'=p'/a_0^2-\rho_0 s'/c_p \tag{3b}$$

$$\frac{D_0}{Dt}\left(\frac{1}{a_0^2}\frac{D_0G}{Dt}\right)-\frac{1}{\rho_0}\nabla\bullet(\rho_0\nabla G)=\frac{1}{\rho_0}\nabla\bullet(\rho_0 v) \tag{4}$$

Here u', s', p', and ρ' are the fluctuating components of velocity, entropy, pressure, and density, while $U=(U_i)$, ρ_0, and a_0 are the velocity, density, and speed of sound for the mean flow. c_p is the specific heat at constant pressure (assumed to be constant). $D_0/Dt=\partial/\partial t+U\bullet\nabla$ is the convective derivative relative to the mean potential flow, and the components of the vector $(X-iU_\infty t)$ are essentially the Lagrangian coordinates of a point convected by the mean flow. Far upstream the mean flow is assumed to be uniform and parallel to the x_1 direction, whose unit vector is i, with speed U_∞. The arbitrary functions b and a with argument $(X-iU_\infty t)$ are chosen to satisfy the upstream boundary conditions for the fluctuations in entropy and vortical velocity. Roughly speaking, $v=(v_i)$ is the velocity associated with the vorticity or entropy fluctuations. Note that Eq. (1) does not correspond to the classical decomposition of a vector field, i.e., the condition $\nabla\bullet v=0$ is not satisfied.

Since far upstream the flow is uniform and parallel, the arbitrary functions b and a generally can be obtained by inspection. The vector function X depends only on the mean flow quantities, and can be determined exactly. Then the only unknown quantity in the above equations is the perturbation potential, G, which satisfies Eq. (4)—an equation with variable coefficients. The source term vanishes in regions where the mean flow is uniform. When rigid surfaces are present, the appropriate boundary condition on the body is

$$n\bullet\nabla G=-n\bullet v \tag{5}$$

where n is the unit vector normal to the surface. The boundary condition as $X\to\infty$ is determined by applying the Sommerfeld radiation or outgoing wave condition.

Now consider a two-dimensional irrotational mean flow, i.e., $U_i=U_i(x_1,x_2)$, $i=1,2$, and $U_3=0$. We will introduce the velocity potential ϕ, and the stream function ψ, of this flow as orthogonal coordinates in the (x_1,x_2) plane. The appropriate coordinate metrics are $h_\phi=(U_0)^{-1}$ and $h_\psi=(\rho_0U_0/\rho_s)^{-1}$, where U_0 is the magnitude of U and ρ_s is an arbitrary constant which will be specified later. For convenience, we will also introduce $z=(U_\infty x_3)$. The convective derivative then becomes $D_0/Dt=(\partial/\partial t+U_0^2\partial/\partial\phi)$.

Noting that our equations are linear, we can represent an arbitrary incoming vortical or entropic disturbance as a superposition of harmonic waves. Far upstream the disturbances should be of the form appropriate for a uniform mean flow. Thus, for the vortical velocity and entropy, we consider an upstream disturbance field of the type

$$v/U_\infty=Ae^{i\alpha},\quad s/2c_p=Be^{i\alpha}\quad i=(-1)^{1/2} \tag{6a,b}$$

where

$$\alpha=(k_t\phi+k_n\psi+k_zz-k_tU_\infty^2t) \tag{6c}$$

Here $k=(k_t,k_n,k_z)$ is the wavenumber and $A=(A_t,A_n,A_z)$ and B are the amplitudes of the vortical velocity and entropy fluctuations. They are arbitrary constants except for the requirement that $A_tk_t+A_nk_n\rho_\infty/\rho_s+A_zk_z=0$, since far upstream the vortical velocity must be divergence free.

We will now develop explicit expressions for the components of rotational velocity, $v=(v_t,v_n,v_z)$. This could be done by writing Eq. (2b) in the curvilinear coordinate system. However, it is simpler to develop the solutions directly in the (ϕ,ψ,z) coordinate system. This also eliminates a difficulty which occurs in Ref. 2 for mean flows past a two-dimensional lifting body.

The second term on the right-hand side of Eq. (2b), which we temporarily call v^*, is the solution of the homogeneous equation

$$D_0v^*/Dt+v^*\bullet\nabla U=0$$

Separating this vector equation into its three components, integrating the resultant first order partial differential equations, enforcing conditions (6) upstream, and substituting the results into Eq. (2b), we obtain the general solution for the vortical velocity

$$v(\phi,\psi,z,t)=(v_t,v_n,v_z)\,v_t/U_\infty=[A_t^*U_\infty/U_0+BU_0/U_\infty]e^{i\alpha'} \tag{7a}$$

$$v_n/U_\infty=(\rho_0U_0/\rho_\infty U_\infty)[A_n+(\rho_\infty A_t^*/\rho_s)\partial g/\partial\psi]e^{i\alpha'} \tag{7b}$$

$$v_z/U_\infty=A_ze^{i\alpha'} \tag{7c}$$

where

$$A_t^*=(A_t-B) \tag{7d}$$

$$\alpha'=k_t\phi+k_tg(\phi,\psi)+k_n\psi+k_zz-k_tU_\infty^2t \tag{7e}$$

$$g(\phi,\psi)=\int_{-\infty}^{\phi}[U_\infty^2/U_0^2(\zeta,\psi)-1]d\zeta \tag{7f}$$

The function $g(\phi,\psi)$ is related to the distortion of vortex filaments by the nonuniform mean flow and is the analog of Lighthill's drift function[4] in the (ϕ,ψ) plane.

We will now transform Eq. (4) for the perturbation potential G into the curvilinear coordinate system. Since the base flow and the boundaries are independent of the spanwise coordinate and time, we can factor out the z and t dependence. Thus setting

$$G=\hat{G}\exp[i(k_zz-k_tU_\infty^2t)] \tag{8a}$$

$$v=\hat{v}\exp[i(k_zz-k_tU_\infty^2t)] \tag{8b}$$

we obtain

$$\frac{\partial}{\partial\phi}\left(\beta_0^2\frac{\partial\hat{G}}{\partial\phi}\right)+\frac{\partial}{\partial\psi}\left(\frac{\rho_0^2}{\rho_s^2}\frac{\partial\hat{G}}{\partial\psi}\right)+2ik_t\frac{U_\infty^2}{a_0^2}\frac{\partial\hat{G}}{\partial\phi}$$

$$+\frac{U_\infty^2}{a_0^2}\left[\frac{k_t^2U_\infty^2}{U_0^2}-\frac{k_z^2}{M_0^2}-2ik_t\frac{\partial}{\partial\phi}(\log a_0)\right]\hat{G}$$

$$=-\frac{\partial}{\partial\phi}\left(\frac{U_\infty}{U_0}\hat{v}_t\right)-\frac{\partial}{\partial\psi}\left(\frac{\rho_0U_\infty}{\rho_sU_0}\hat{v}_n\right)-ik_z\frac{U_\infty^2}{U_0^2}\hat{v}_z \tag{8c}$$

Here $M_0=U_0/a_0$ is the local Mach number of the mean flow and $\beta_0^2=(1-M_0^2)$. Although Eq. (8c) is linear, it has variable coefficients and first order derivatives are present. Thus, except for special cases, it appears that solution by numerical methods would be necessary.

It will now be shown that, when the thermodynamic behavior of the mean flow is assumed to follow the tangent gas relations,[3] Eq. (8c) can be simplified significantly. The tangent gas approximation consists of a linearization of the equation of state in the pressure-specific volume plane. For our application we choose the tangency point about which the linearization is made to be the mean flow conditions at infinity. We then obtain

$$a_0\rho_0 = a_\infty\rho_\infty \qquad a_0\beta_0 = a_\infty\beta_\infty \tag{9a,b}$$

Introducing the transformed velocity potential, which is a generalization of the Miles transformation,

$$H = \hat{G}\beta_0 \exp[ik_t M_\infty^2 \phi/\beta_\infty^2] \tag{10a}$$

utilizing Eqs. (9) and choosing the arbitrary constant $\rho_s = \rho_\infty/\beta_\infty$ we can write Eq. (8c) in the form

$$\frac{\partial^2 H}{\partial\phi^2} + \frac{\partial^2 H}{\partial\psi^2} + \left[\left(k_t^2\frac{M_\infty^2}{\beta_\infty^2} - k_z^2\right)\left(\frac{M_\infty^2}{\beta_\infty^2} + \frac{U_\infty^2}{U_0^2}\right) - \frac{1}{\beta_0}\left(\frac{\partial^2\beta_0}{\partial\phi^2} + \frac{\partial^2\beta_0}{\partial\psi^2}\right)\right]H = S(\phi,\psi) \tag{10b}$$

where

$$S(\phi,\psi) = -ie^{i\gamma}\left\{\frac{k_t A_t^*}{\beta_0}\left[\frac{U_\infty^4}{U_0^4} - \frac{U_\infty^2}{U_0^2} + \left(\beta_0\frac{\partial g}{\partial\psi}\right)^2\right] + k_n A_n\left(\frac{\beta_0}{\beta_\infty} - \frac{U_\infty^2\beta_\infty}{U_0^2\beta_0}\right) + \left(k_n A_t^* + \frac{k_t A_n}{\beta_\infty}\right)\beta_0\frac{\partial g}{\partial\psi}\right\}$$
$$- e^{i\gamma}\left\{\frac{A_t^*}{\beta_0}\left[\frac{\partial}{\partial\psi}\left(\beta_0^2\frac{\partial g}{\partial\psi}\right) - 2\frac{U_\infty^2}{U_0^3}\frac{\partial U_0}{\partial\phi}\right] + 2\frac{A_n}{\beta_\infty}\frac{\partial\beta_0}{\partial\psi}\right\} \tag{10c}$$

and

$$\gamma = k_t\phi/\beta_\infty^2 + k_t g(\phi,\psi) + k_n\psi \tag{10d}$$

The transformed boundary condition [Eq. (5)] for any rigid surfaces which may be present is

$$\left(\frac{\partial H}{\partial\psi} - \frac{1}{\beta_0}\frac{\partial\beta_0}{\partial\psi}H\right)_{\psi=\psi_0} = -\frac{\beta_0}{\beta_\infty}\left[A_n + \beta_\infty A_t^*\frac{\partial g}{\partial\psi}\right]e^{i\gamma} \tag{10e}$$

where $\psi = \psi_0$ on the body surface and the range of ϕ corresponds to the body length in the (ϕ,ψ) plane. Equation (10) are the key results of this Note.

Equation (10b) for the transformed velocity potential has several advantages over Eq. (8c) or Eq. (4). The transformation has eliminated the first-order derivatives of H and the only variable coefficient in the differential operator is that multiplying H. This is a significant simplification as compared to Eq. (8c) or Eq. (4), which have first-order derivatives and variable coefficients in the derivative terms. The boundary condition on any rigid surface is also applied along a coordinate line $\psi_0 = \text{const}$, which is not the case in general for Eq. (4).

The form of Eq. (10b) is also very useful in that the qualitative character of the solution is easily recognized. For positive values of the coefficient of H the solution will exhibit propagating wave behavior, while for negative values the solution will decay exponentially. The propagation wave behavior is generally of the most interest. We see that the term multiplying $(k_t^2 M_\infty^2/\beta_\infty^2 - k_z^2)$ in the coefficient of H is always positive. It will also be shown below that $(\nabla^2\beta_0)/\beta_0$ is always negative. Thus, if $(k_t^2 M_\infty^2/\beta_\infty^2 - k_z^2) > 0$, the solution for the modified velocity potential H will exhibit propagating wave behavior over the entire (ϕ,ψ) plane. Under these conditions, Eq. (10b) may be thought of as an inhomogeneous Helmholtz equation describing wave propagation in a medium with variable "index of refraction."

We will now discuss the second term of the coefficient of H in Eq. (10b). For the tangent gas theory,[3] it can be shown that

$$\beta_0 = \tanh(\Omega + \Omega^*)$$

where

$$\Omega(\phi,\psi) = \int_{U_\infty}^{U_0} (\beta/U)\,dU$$

is a speed variable which satisfies Laplaces equation and Ω^* is a constant (i.e., independent of ϕ and ψ). Thus we obtain

$$\frac{\nabla^2\beta_0}{\beta_0} = \frac{1}{\beta_0}\left(\frac{\partial^2\beta_0}{\partial\phi^2} + \frac{\partial^2\beta_0}{\partial\psi^2}\right) = -2\,\text{sech}^2(\Omega+\Omega^*)\left[\left(\frac{\partial\Omega}{\partial\phi}\right)^2 + \left(\frac{\partial\Omega}{\partial\psi}\right)^2\right]$$

which is always less than or equal to zero. We also note that in many applications the term $(\nabla^2\beta_0)/\beta_0$ can be neglected without any appreciable error. One case is that in which the disturbance wavelength is much smaller than the scale on which the mean flow varies. The first term in the coefficient of H will then dominate and the solution will have a geometric acoustics behavior. A second case is that of mean flows which are a small disturbance (say $0(\epsilon)$) to a uniform flow. The term $(\nabla^2\beta_0)/\beta_0$ is then $0(\epsilon^2)$ and thus negligible in comparison to the $0(\epsilon)$ first term in the coefficient of H.

For mean flows which are a small disturbance to a uniform stream, the source term in Eq. (10b) can also be simplified significantly. The speed variable $\Omega = 0(\epsilon)$, and we have

$$U_\infty/U_0 \cong 1 + \Omega/\beta_\infty \qquad \beta_0 \cong \beta_\infty + M_\infty^2\Omega \tag{11a,b}$$

where terms of $0(\epsilon^2)$ have been neglected. To $0(\epsilon)$, Eq. (7f) for the drift function then involves an integral of the form $\int\Omega(\phi,\psi)\,d\phi$. The straightforward evaluation of this integral is often difficult, since the expressions for Ω are generally quite complicated. However, we can take advantage of the fact that Ω is the real part of an analytic function[3] $\tau(w) = \Omega(\phi,\psi) + i\theta(\phi,\psi)$ where θ is the angle of the local velocity with respect to the uniform velocity at infinity and $w = \phi + i\psi$. Thus we obtain

$$\int\Omega(\phi,\psi)\,d\phi = \text{Re}[T(w)] \qquad T(w) = \int\tau(w)\,dw \tag{12a,b}$$

The complex integral (12b) is independent of path (as long as the standard restrictions are met) and is generally much easier to evaluate than $\int\Omega\,d\phi$. We then obtain for the drift function

$$g(\phi,\psi) \cong \frac{2}{\beta_\infty}\,\text{Re}[T(w)] \qquad \lim_{R\to\infty}\{\text{Re}[T(Re^{i\pi})]\} = 0 \tag{13a,b}$$

where the arbitrary constant in $T(w)$ is chosen in accordance with Eq. (13b). Setting

$$d\tau/dw = \Omega'(\phi,\psi) + i\theta'(\phi,\psi) \tag{14a}$$

the source term given by Eq. (10c) simplifies to

$$S(\phi,\psi) \cong -2e^{i\gamma}\{i(k_t A_t^*/\beta_\infty^2 - k_n A_n \beta_\infty)\Omega$$

$$-i(k_n A_t^* + k_t A_n/\beta_\infty)\theta$$

$$+(A_t^* M_\infty^2/\beta_\infty^2)\Omega' - (A_n M_\infty^2/\beta_\infty)\theta'\} \qquad (14b)$$

where terms of $0(\epsilon^2)$ have been neglected.

It is also useful to note that $T(w)$ is related to the complex velocity potential for an incompressible flow which is a transformation of the mean flow past the body of interest. The perturbation velocity potential for the linearized compressible mean flow satisfies $\beta_\infty^2 \partial^2\Phi/\partial x^2 + \partial^2\Phi/\partial y^2 = 0$ and appropriate conditions on the body surface. Here (x,y) are Cartesian coordinates aligned with the uniform flow at infinity. Expanding the flow speed and angle in terms of the perturbation velocities we find to linear order

$$\partial\Phi/\partial x = -U_\infty \Omega/\beta_\infty \quad \partial\Phi/\partial y = U_\infty \theta \qquad (15a,b)$$

We now introduce the coordinate transformation

$$\Phi'(x',y') = \Phi(x,y) \quad x' = x, y' = \beta_\infty y$$

The function Φ' then satisfies Laplace's equation and is the real part of an analytic function $F(z') = \Phi'(x',y') + i\Psi'(x',y')$ of $z' = x' + iy'$. $F(z')$ is determined by using the transformed boundary conditions on the body and $dF/dz' \to 0$ as $z' \to \infty$. Applying the coordinate transformation to Eqs. (15) and noting that $w = U_\infty z' + 0(\epsilon)$, we find

$$T(w) \cong -\beta_\infty F(w/U_\infty) + \text{const} \qquad (16)$$

where terms of $0(\epsilon^2)$ have been neglected. The constant in Eq. (16) is determined by applying Eq. (13b).

In summary, we have shown that the equations governing disturbances to a subsonic compressible flow can be simplified significantly by using the tangent gas approximation for the mean flow. The transformed equations show that, for $(k_t^2 M_\infty^2/\beta_\infty^2 - k_z^2) > 0$, the irrotational portion of the disturbance will exhibit propagating wave behavior throughout the flowfield. When the mean flow perturbation is small, the equations can be further simplified. All quantities in the equations are then given explicitly in terms of a complex velocity potential and its derivatives. The complex velocity potential can be determined by a simple transformation of the equation for the linearized compressible mean flow past the body of interest.

Acknowledgment

The work described herein was financially supported by the General Electric Company Aircraft Engine Business Group under the Independent Research and Development Program.

References

[1] Hunt, J.C.R., "A Theory of Turbulent Flow Round Two-Dimensional Bluff Bodies," *Journal of Fluid Mechanics*, Vol. 61, 1973, pp. 625-706.

[2] Goldstein, M. E., "Unsteady Vortical and Entropic Distortions of Potential Flows Round Arbitrary Obstacles," *Journal of Fluid Mechanics*, Vol. 89, 1978, pp. 433-468.

[3] Woods, L. C., *The Theory of Subsonic Plane Flow*, Cambridge University Press, London, 1961.

[4] Lighthill, M. J., "Drift," *Journal of Fluid Mechanics*, Vol. 1, 1956, pp. 31-53.

AIAA 81-4278

Forebody Drag Reduction

B. N. Pamadi*

Indian Institute of Technology, Bombay, India

I. Introduction

A BLUNT body traveling at high speeds through the atmosphere experiences aerodynamic heating. One of the ways to alleviate this effect is to eject a cool gas jet directed against the air stream from the nose of the body.[1-3] The ejected gas flows back over the nose of the body and forms a protecting blanket of low temperature fluid. When the momentum of the jet efflux exceeds a certain critical value, the bow shock wave bulges out and stands away from the body surface and takes a form appropriate to a new body consisting of the original body with a protrusion due to jet flow (Fig. 1). The boundary of this protrusion is called the interface, the stream surface between the jet flow and the main stream flow. A conical, "dead air region" is formed around the jet and in the vicinity of the nose. The pressure on the body within the dead-air region is greatly reduced compared with pressures for no ejection (Fig. 1b). Thus, in addition to heat flux reductions, the forebody pressure drag also decreases considerably. Warren[1] considered this problem mainly with the objective of finding ejection conditions for achieving optimum heat flux reductions. Romeo and Sterret[2] studied the effect of a blowing gas jet on the form and location of bow shock wave. Finley[3] concerned himself mainly with the task of semiempirical prediction of the flow model. Both Warren[1] and Finley[3] have reported several pressure measurements over the forebody. However, full note of the large reductions in the forebody pressure drag has not been taken by these authors. It also appears that this point has not received further attention in literature.

In this Note the forebody pressure distributions taken from the studies of Warren[1] and Finley[3] have been integrated to deduce the corresponding drag coefficients. From this exercise, it is observed that significant forebody drag reductions can be achieved by this technique. Further, it is also shown that these data can be effectively correlated using a momentum coefficient which characterizes the jet efflux and freestream conditions.

II. Description of the Flow Pattern

We introduce a momentum coefficient to characterize the jet efflux

$$C_F = \left[p_j \frac{\pi}{4} d_j^2 + w_j v_j\right] \Big/ \left[\frac{1}{2}\rho_\infty V_\infty^2 \frac{\pi}{4} d_m^2\right] \qquad (1)$$

where p_j is the static pressure at jet exit, v_j the jet exit velocity, w_j the jet mass flow rate, and V_∞ the freestream velocity.

For small values of momentum of the jet efflux ($C_F \leq 0.015$), no noticeable change occurs in the basic flow pattern, which comprises of a strong bow shock wave standing close to the body. When C_F exceeds 0.015, the bow shock wave begins to bulge out. Except for a small range of total pressures of the jet when the flow pattern is unsteady, the

Received Oct. 16, 1980; revision received May 1, 1981.

*Assistant Professor, Dept. of Aeronautical Engineering; presently, NRC-NASA Associate, NASA Langley Research Center, Hampton, Va..