van den Dungen's method must depend on the small difference between two large and somewhat uncertain terms, whereas in our case the difference is of the same magnitude as the terms from which it is formed.

2. The pressure data in Table 2 go back largely to Napier Shaw, the source used also by van den Dungen in his computations. The fact that the pressure changes do not integrate to zero over the whole globe can be ascribed to uncertainties in these observations, and in the conversion of sea level pressure to local pressure. We have accordingly stated that "it would be desirable to recompute Table 2 using the original data for the uncorrected local pressure". In any event it can be shown that our estimate for the magnitude of the inertia term (which is small compared to the wind term) will hardly be effected by allowing for this discrepancy.

3. We regret to have given the impression that van den Dungen et al were unaware of the importance of reverting the sea level pressures to local pressures. The complete corrections involve the *elevation* and the *temperature* of the weather stations. The corrections, as entered in Table 1 of van den Dungen's paper, take into account only the elevations. This part of the correction, which is of course the same throughout the year, will drop out from computations of the seasonal *differences* with which our problem is concerned. This explains why van den Dungen's February minus August values (their Table 4) are nearly the same as our *uncorrected* January minus July values (our Table 2, column 2).

The complete corrections depend partly on *temperature*, and these corrections do not drop out when differences between seasons are computed. The corrections are remarkably large. North of 30° N latitude the average local pressure variation between January and July is about *minus* 1.5 mb, whereas van den Dungen et al. obtained roughly *plus* 4 mb.

Comments on the Behavior of Jet Streams over Eastern North America

By P. F. CLAPP and J. S. WINSTON
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October 11th, 1950

Mr Phillips' article in the May 1950 issue of TELLUS (Vol. 2, No. 2) touches on an interesting point regarding confluence. It is quite true that the mere confluence of streamlines on upper level charts does not necessarily lead to increased thermal gradients and consequent increases in high level wind speeds. Confluence of the streamflow must be accompanied by differential temperature advection at most levels of the troposphere before it can result in accelerating wind speeds in the upper troposphere. However, Mr Phillips fails to consider that lack of strong advection at 300 mb may not represent conditions at all levels. It is very likely that there is cold air advection in the northwesterly flow and warm air advection in the southwesterly flow at middle and lower levels of the troposphere (i.e. up to about 400 mb). Hydrostatically, the mean temperature field of the layer below 300 mb must be similar in appearance to the 300 mb height field, but it is well-known that the flow at lower levels in middle latitudes is generally less zonal and more meridional than at 300 mb. Thus, even in the case Mr Phillips presents the streamline-isotherm relationship prescribed by the confluence theory of NAMIAS and CLAPP (J. Meteor., 1949, Vol. 6, pp. 330—336) probably does exist at most levels of the troposphere. In fact examination of the pertinent sea level and 700 mb charts indicates that this is so in the layer from sea level to 10,000 feet. Our opinion, contrary to Mr Phillips' implication, is that the situation described by Namias and Clapp is the frequent rather than the rare occurrence in cases of confluence.

Note on recent study of stability by R. Fjørtoft

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August 1950

The purpose of this note is to call attention to a method for investigating the dynamic stability of certain flowsystems recently presented by R. FJØRTOFT (1950), and to show that his results may be expressed in a clear and concise form through the use of the quasi-Lagrangian system of hydrodynamical equations discussed by the writer (STARR 1945) on a previous occasion. In order to achieve clarity, the line of reasoning employed by Fjørtoft is here repeated in revised form.

Let us consider the two-dimensional horizontal motion of an inviscid incompressible fluid of uniform density on the surface of a sphere of radius r, in the region between two fixed latitudes Φ_1, and Φ_2 at which walls are assumed to be present so that a complete zonal belt extending through all longitudes λ is formed. Following Fjørtoft we examine the area integral I of the absolute vorticity ζ multiplied by $\sin \Phi$ over the entire area A of the belt. Either by introducing the proper definition of ζ in terms of the particle velocity and integrating, or by the following application of Green's theorem, I can be shown to remain invariant for arbitrary motions of the fluid. Since a stream function Ψ must exist such that $ru = -\partial\Psi/\partial\Phi$; $rv \cos \Phi = \partial\Psi/\partial\lambda$; $\zeta = \nabla^2\Psi$, where u and v are the eastward and northward components of particle velocity and ∇^2 is the Laplacian operator for the curved surface, we have that

$$I \equiv \int\int \sin \Phi \cdot \nabla^2 \Psi dA = -\oint \sin \Phi \cdot \frac{\partial \Psi}{\partial n} ds - \int\int r^{-2} \cos \Phi \cdot \frac{\partial \Psi}{\partial \Phi} dA.$$

Here ds is an element of the boundary of A and n is the outward normal to it. Because the region is doubly-connected, the line integral involves the use of a cut, but due to the fact that the integrand is single-valued the cut may be eliminated and the integral separated into two cyclic line integrals along Φ_1, and Φ_2. Upon substituting the particle velocities for the derivatives of Ψ it may be observed that I is equal to a linear combination with constant coefficients of (a) the circulation around Φ_1, (b) the circulation around Φ_2, and (c) the total angular momentum of the fluid (contained in the last integral). In view of the fact that the fluid chains next to the walls must remain there as the motion progresses and that no solenoids or friction are present, each of these three quantities and therefore also I must remain constant with time t.

Examining I otherwise, we may first change the variables of integration from Φ, λ to y, λ where $y = \sin \Phi$. Furthermore we may then in turn change these variables from y, λ to y_0, λ where y_0 is the average value of y with respect to longitude for a given material chain of particles encircling the sphere, according to the quasi-Langrangian scheme of hydrodynamical equations discussed elsewhere (STARR 1945). Thus

$$I \equiv r^2 \int\int \sin \Phi \cdot \zeta \cdot \cos \Phi d\Phi d\lambda = r^2 \int\int y\zeta dy d\lambda = r^2 \int\int y\zeta \frac{\partial y}{\partial y_0} dy_0 d\lambda.$$

If this is applied to motion which represents a breakdown from purely zonal flow, we may choose the material chains in such a fashion that ζ, $\frac{d\zeta}{dy_0}$ are functions of y_0 alone. In view of the fact that under the circumstances assumed $d\zeta/dt \equiv 0$, this functional dependence on y_0 is then preserved in the subsequent development of the motion. Supposing further that we define ε by the relationship $y = y_0 + \varepsilon$, it is simple to show that from continuity considerations $\oint \varepsilon d\lambda = 0$ for each chain $y_0 =$ constant. Upon making use of these several conditions it follows that

$$I = r^2 \int\int (y_0 + \varepsilon) \zeta \left(1 + \frac{\partial \varepsilon}{\partial y_0}\right) dy_0 d\lambda = r^2 \int\int y_0 \zeta dy_0 d\lambda - \frac{r^2}{2} \triangle,$$

where the first integral on the right is the value of I corresponding to zonal flow, i.e., when $\varepsilon \equiv 0$, and $y_0 \equiv y$, and $\triangle$ is given by

$$\triangle \equiv \int \left[\frac{d\zeta}{dy_0} \oint \varepsilon^2 d\lambda\right] dy_0.$$

Due to the invariant character of I demonstrated before, it is clear that $\triangle$ must vanish for all motions developing from a state of zonal flow and for all motions which develop into zonal flow as time progresses. Several cases may be distinguished.

(1) If $d\zeta/dy_0 > 0$ everywhere from $y_0 = y_1$, to $y_0 = y_2$ or if $d\zeta/dy_0 < 0$ everywhere in the same interval of y_0, then $\triangle$ cannot vanish unless $\varepsilon \equiv 0$. In these cases zonal flow cannot develop into disturbed motion.

(2) If $d\zeta/dy_0$ changes sign or becomes zero within the interval $y_0 = y_1$ to $y_0 = y_2$, we *may* have $\triangle = 0$ and $\varepsilon \neq 0$. In this case zonal flow may develop into disturbed motion or vice versa, but no information is given by the present analysis as to whether it will actually do so.

(3) Preserving the conditions given under (1) let us suppose that a disturbance is introduced forcibly by some unspecified external agency, after which the system is left to its own devices. It follows that the finite value of $\triangle$ so introduced cannot of its own accord disappear, so that the motion will then continue to be disturbed.

It is interesting to note the following. Under (1) stability is implied whatever be the distribution of ζ itself so long as the derivative of ζ has the properties stated. This indicates that criteria which relate the presence or absence of stability to the value of ζ or to the gradient of angular momentum are not applicable to problems of the present type where viscosity is absent and the motions are two-dimensional. On the other hand, the results recapitulated in this note are in agreement with those obtained by LIN (1945) and KUO (1949).

For certain disturbances covered under (3), according to LIN (1945) and KUO (1949), the introduction of a very small viscosity is effective in producing damping without at the same time leading to a sensible dissipation of the total kinetic energy. The introduction of such viscosity cannot, however, produce instability for the cases covered under (1).

The invariance of I is in essence a spherical adaptation of the principle of invariance of the centroid of a vortex system as demonstrated, for example, in the textbook by LAMB (1932). The extremal properties of I (as here demonstrated by the character of $\triangle$) for certain circumstances, are mathematically similar to those of the geopotential energy integral for properly selected stratified fluids as pointed out by FJØRTOFT (1950).

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Floating ice islands in the Arctic Ocean

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During his recent visit to Alaska, Dr Sverre Petterssen suggested that the readers of Tellus might be interested in information on the recent discoveries of large floating ice islands in the Arctic Ocean.

The undersigned, having reason to suspect the existence of such islands, initiated visual and radar search in Spring 1950 by the 375th Reconnaissance Squadron in connection with the routine weather reconnaissance flights to the North Pole (the Ptarmigan Flight). Although no deviation from the standard track could be made in search of such islands, the suspicion was soon corroborated. On June 3, 1950 the radarscope showed indications of a large floating island (called T-1) at 73°15'N and 159°05'W. This, or a similar ice mass was seen by radar ten days later at 74°30'N and 159'05'W. Whether or not these radar returns were from the same or from two different ice masses could not be determined. The relatively large distance between them might, however, suggest that they were two different islands.

Fig. 1. At a distance of 65 miles, T-2 could be distinguished clearly from the surrounding pack. This photograph was taken at a distance of 45 miles.

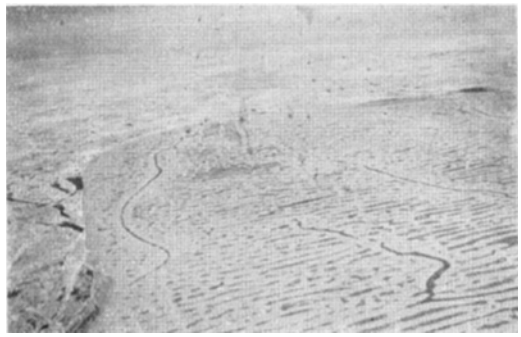

Fig. 2. The high edge of T-2. Note the washboard pattern, the drainage system and the steep cliff.

The first visual identification of floating islands was made at 86°40'N and 167°00'E on 20 July 1950 when three photographs were taken; during the following two weeks several additional pictures were obtained. This island has been called T-2.

Island T-3 was discovered toward the end of July 1950 when several radar oscilloscope photographs were taken of a rather small but well-defined island at approximately 75°24'N and 173°00'W.