

A Simple Theory for Waterspouts

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ABSTRACT

It is shown that the simple thermodynamic theory for dust devils, proposed by Rennó et al., also applies to waterspouts. The theory is based on the thermodynamics of heat engines and predicts the central pressure and the wind speed of these convective vortices. Moreover, it provides a simple physical interpretation of their general characteristics. In particular, the heat engine theory shows that convective vortices are more likely to form in the regions where the occurrence of the warmest and moistest updrafts and the coldest and driest downdrafts are supported by the local environment. These are the regions where both the heat input into the convective heat engine is maximum and the solenoidal generation of vorticity is the greatest. This explains why waterspouts are frequently observed near the boundaries between relatively warm and relatively cold waters. Moreover, since the work done by the convective heat engine is equal to the total heat input multiplied by the thermodynamic efficiency, the theory shows that another necessary condition for the formation of intense vortices is the presence of intense convection.

1. Introduction

The beauty of dust devils and waterspouts stimulates the imagination of the layman or scientist who observes one of them. Indeed, these vortices are not only of exceptional natural beauty but also of great scientific interest. They are of particular meteorological interest because they might be thermodynamically similar to hurricanes or tornadoes, the most destructive storms observed on earth. Therefore, the theoretical understanding of these smaller and weaker convective vortices has the potential to further our understanding of their larger and violent cousins. Since dust devils and waterspouts are relatively frequent, small, and weak, they allow safer and more efficient *in situ* measurements. Convective vortices are all low-pressure warm-core vortices with radius of maximum wind varying between about 1 m (dust devils) and 25 km (hurricanes). In this article, we focus on the study of waterspouts, which cover a wide range of sizes and intensities. Thorough reviews on the observations of dust devils are given by Sinclair (1966, 1969, 1973).

Golden (1968, 1971, 1973, 1974a,b, 1977), Levenson et al. (1977), Simpson et al. (1986), and Wakimoto and

Lew (1993) describe the characteristics of waterspouts in detail. Here we briefly review some of their findings. Hundreds of waterspouts are observed each summer around the Florida Keys, and a small number over the open oceans. Waterspouts have surface diameters between 5 and 75 m. Like their smaller cousins, dust devils, they receive their vorticity from local wind shear. Therefore, they rotate either cyclonically or anticyclonically (Golden 1974a; Schwiesow 1981). Figure 1 is a sketch of a waterspout. We speculate that, near the surface, air parcels absorb heat from the surface as they spiral in toward the waterspout and become warmer and moister than the ambient air. The temperature and pressure perturbations observed within waterspouts vary from 0.2 to 2.5 K and from 10 to 90 hPa (Golden 1974a; Levenson et al. 1977). The vertical velocity reaches positive peak values of about 10 m s^{-1} in the region of highest temperature. In the waterspout there is solid-body rotation and a weak forced downdraft in its core (Golden 1974a). A more intense rain-cooled downdraft is frequently observed near the waterspout. The waterspout's low-level tangential velocity can reach peak values of about 80 m s^{-1} (Golden 1974a; Schwiesow 1981).

To a first approximation, the tangential winds of waterspouts are in cyclostrophic balance. However, since near the surface there is a radial inflow of air toward their center, the observed pressure gradients are slightly larger than those necessary to support the cyclostrophic

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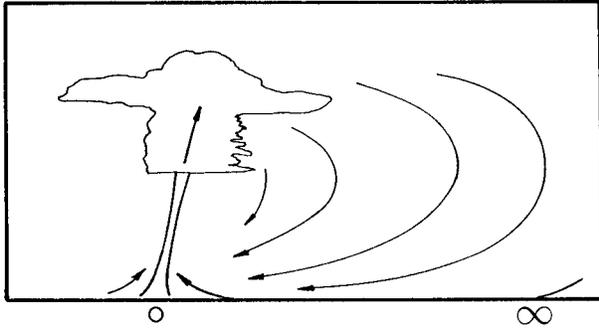


FIG. 1. Sketch of a waterspout (not in scale) indicating the near-surface integration path [from large radius (∞) to the center of the vortex (\circ)] along the inflow streamtube and the material volume occupied by the convective system. The inflow streamtube is defined by the interception of the streamlines with the surface layer (the dotted area at the bottom of the figure).

tangential winds. Waterspouts have tangential velocity profiles characteristic of a Rankine vortex. That is, near their center the tangential velocity is proportional to the radius; from there it decreases inversely proportional to it.

Waterspouts usually form under convective clouds deeper than about 3 km (Golden 1974a). They usually originate in a region of horizontal wind shear between updrafts and downdrafts (Golden 1974a; Hess and Spillane 1990). Moreover, the observations suggest that the existence of these local horizontal shear lines (vertical vorticity) separating the updrafts from the downdrafts is a necessary, but not sufficient condition for waterspout formation (Simpson et al. 1986). It is important to point out that these shear lines can be generated by the convective system itself (see Simpson et al. 1991).

Our main objective is to propose a simple scaling model for the potential intensity of waterspouts. In order for a convective vortex to form, both thermodynamic processes responsible for maintaining a pressure depression and dynamic processes capable of producing vorticity must be present. The various dynamic processes capable of producing and enhancing vertical vorticity in convective systems have been extensively studied in the atmospheric science literature (Lilly 1982; Davies-Jones 1984; Rotunno and Klemp 1985; Simpson et al. 1986). We do not discuss these mechanisms in this paper. Instead, we focus on the thermodynamics of the convective process responsible for the maintenance of the pressure depression within a convective vortex. We consider a waterspout vortex that is in quasi-steady state and cyclostrophic balance. Since for a given set of environmental conditions, a steady state is achieved only when the work done by the heat engine is balanced by mechanical friction, the quasi-steady-state assumption implies that we aim at a simple theory for the maximum bulk thermodynamic intensity of a vortex in cyclostrophic balance.

2. Theoretical framework

Emanuel (1986), Emanuel (1989), Rennó and Ingersoll (1996), Emanuel and Bister (1996), Michaud (1977, 1995), and Rennó et al. (1998) idealized hurricanes, polar lows, convective systems, and dust devils as heat engines. In this study, we show evidence that the framework proposed by Rennó et al. (1998), for the understanding of dust devils, also applies to waterspouts. As in the study of Rennó et al. (1998), this study focuses on the understanding of waterspout vortices from the standpoint of the thermodynamics, rather than from the standpoint of the mechanisms that generate vorticity. Our goal is to provide a scaling theory for the potential intensity of convective vortices as a function of thermodynamic variables. Naturally, the presence of convection alone is not sufficient for the genesis of a convective vortex. Mechanisms for the generation of vorticity are also necessary. However, since the flow around dust devils, and all but perhaps the most intense waterspouts are, to a first approximation, in cyclostrophic balance, their wind speed is determined by the bulk pressure drop from the ambient to their center. The bulk pressure drop, in turn, depends solely on the thermodynamics of their convective heat engine, as explained below. Thus, to a first approximation, the wind speed around these vortices is determined by the thermodynamics of their convective heat engine.

An atmospheric vortex in cyclostrophic balance satisfies the equation

$$\frac{v_a^2}{a} \approx \alpha \left(\frac{\Delta p}{a} \right), \quad (1)$$

where a is the vortex radius (the radius of maximum wind), v_a is the tangential wind speed at a , α is the specific volume of moist air, and Δp is the radial pressure drop across the vortex.

It follows from Eq. (1) and the ideal gas law that, to a first approximation, the tangential wind speed around a convective vortex does not depend on its radius; that is,

$$v_a \approx \sqrt{\frac{RT_\infty \Delta p}{p_\infty}}, \quad (2)$$

where R is the specific gas constant for moist air, and T_∞ and p_∞ are the surface temperature and pressure in the environment, at the vortex radius of influence.

Rennó et al. (1998) showed that the pressure drop across a convective vortex [see their Eq. (16)] is given by

$$\begin{aligned} \Delta p &\equiv (p_\infty - p_0) \\ &\approx p_\infty \left\{ 1 - \exp \left[\left(\frac{\gamma \eta}{\gamma \eta - 1} \right) \left[\left(\frac{c_p}{R} \right) \left(\frac{T_0 - T_\infty}{T_\infty} \right) \right. \right. \right. \\ &\quad \left. \left. \left. + \left(\frac{L_v}{R} \right) \left(\frac{r_0 - r_\infty}{T_\infty} \right) \right] \right] \right\}, \quad (3) \end{aligned}$$

where p_0 and T_0 are the temperature and pressure at the vortex center, γ is the fraction of the total dissipation of mechanical energy consumed by friction near the surface, η is the thermodynamic efficiency, c_p is the heat capacity of moist air per unit mass, L_v is the latent heat of vaporization of water per unit mass, and r_0 and r_∞ are the water vapor mixing ratio at the vortex center and at large radius (∞). It is important to note that, in the definition of γ , the mechanical energy consumed by friction near the surface is due to both friction between the surface and the near-surface air parcels converging toward the vortex center, and friction throughout the vortex inflow region. Since a steady state is reached only when the frictional loss of energy is balanced by the net heat input, a steady state is possible only in the presence of friction. Therefore, in our theory γ is always larger than zero. Moreover, in the inviscid limit γ is *not* equal to zero. In this case, γ approaches a singular value, that is, $\lim_{\nu \rightarrow 0} \gamma = 0/0$, where ν is the fluid's viscosity. This, in turn, happens because convective vortices cannot reach steady state in inviscid flows.

The following approximations were made in the derivation of Eq. (3). (i) The heat engine cycle is reversible {the transfer of heat from the ground to the near-surface air and the mechanical dissipation of energy are irreversible processes, but both are considered to be external to the heat engine cycle [Rennó and Ingersoll (1996); see also Emanuel (1997) and Bister and Emanuel (1998) for discussions of the effects of dissipative heating]}. (ii) There is a stagnation point at the center of the convective vortex, and to a first approximation the surface is flat. Therefore, changes in kinetic and potential energy of air parcels moving toward the center of the vortex¹

¹ In a waterspout vortex that is two celled (e.g., Golden 1974a), the radial component of the wind vanishes underneath the rising branch of the inner and outer cells. Equation (3) still holds along a streamline at the surface, where the vertical velocity vanishes. However, since most of the pressure deficit in numerically simulated intense vortices is associated with flow at larger radii beyond where the radial wind component vanishes [see Figs. 6 and 10 in Rotunno (1984)], Eq. (3) is a good approximation even when Eq. (1) in Rennó et al. (1998) is integrated radially inward only to the point at which radial wind component vanishes; this point lies well within the radius of maximum azimuthal wind (Rotunno 1984), so that the point at which the radial wind component vanishes, but the azimuthal wind component does not vanish exactly, can be considered approximately as a stagnation point. That is, integrating Eq. (1) of Rennó et al. (1998) inward not all the way to the center of the vortex, but stopping underneath the rising branch of the outer and inner cells, is also a viable procedure.

can be neglected. (iii) The vortex is in steady state, implying that Eq. (3) predicts the maximum thermodynamic intensity of a convective vortex. Moreover, in Eq. (3) we have approximated the average surface temperature and pressure by their values at large radius. It is important to note that Eq. (3) is general, that is, it includes non-Boussinesq effects (e.g., anelastic effects).

The thermodynamic efficiency of a reversible heat engine is defined as

$$\eta \equiv \left(\frac{\int_{\text{in}} T ds - \int_{\text{out}} T ds}{\int_{\text{in}} T ds} \right), \quad (4)$$

where the subscripts "in" and "out" represent, respectively, integration at the heat intake and at the heat rejection (output) branch of the cycle. Assuming that the heat engine is reversible, Eq. (4) can be written as

$$\eta = \left(\frac{T_h - T_c}{T_h} \right), \quad (5)$$

where

$$T_h = \frac{1}{\Delta s} \int_{\text{in}} T ds \quad \text{and} \quad (6)$$

$$T_c = \frac{1}{\Delta s} \int_{\text{out}} T ds, \quad (7)$$

are the entropy-weighted mean temperatures of the layers where heat is absorbed (the hot source) and where heat is rejected (the cold sink), and $\Delta s = \int_{\text{in}} ds = \int_{\text{out}} ds = s_2 - s_1$ is the entropy excess of the convective updrafts over the downdrafts. Thus, for a convective system in steady state, the temperature of the hot source can be approximated by the ambient surface temperature ($T_h \approx T_\infty$), and the temperature of the cold sink (T_c) is equal to the entropy-weighted mean temperature of the convective layer (Rennó and Ingersoll 1996; Rennó et al. 1998). Since, to a first approximation, the mean entropy of a convective layer is constant with height, the entropy-averaged temperature of the convective layer air is equal to its pressure-averaged temperature (Rennó et al. 1998).

Substituting Eq. (3) into Eq. (2), we get an expression for the wind speed around a convective vortex

$$v_a \approx \sqrt{RT_\infty \left\{ 1 - \exp \left[\left(\frac{\gamma \eta}{\gamma \eta - 1} \right) \left[\left(\frac{c_p}{R} \right) \left(\frac{T_0 - T_\infty}{T_\infty} \right) + \left(\frac{L_v}{R} \right) \left(\frac{r_0 - r_\infty}{T_\infty} \right) \right] \right] \right\}}. \quad (8)$$

Equation (8) shows that the maximum intensity of a convective vortex depends only on the thermodynamics of its convective heat engine. That is, it does not depend explicitly on the mechanisms responsible for the generation of vorticity. However, the radius of maximum wind of the convective vortex does depend on the value of the background vorticity. This is shown next.

In the absence of torques acting on the air parcels spiraling in toward the vortex center, their angular momentum is conserved. Thus, the magnitude of their angular momentum per unit mass at a distance d from the vortex's center is

$$l = vd + \frac{1}{2}fd^2 = \text{const}, \tag{9}$$

where f is twice the local vertical component of the earth's angular velocity. In reality, viscous forces reduce the angular momentum of air parcels spiraling in toward the vortex center.

Since for waterspouts the vertical component of the earth's angular velocity at the radius of maximum tangential wind speed can be neglected, the radius of maximum wind speed of a convective vortex is given by

$$a \equiv d_a = \frac{\beta l_\infty}{v_a}, \tag{10}$$

where $(1 - \beta)$ is an air parcel fractional loss of angular momentum as it spirals in from infinity to the vortex center, and l_∞ is the initial angular momentum per unit mass of air parcels spiraling in toward the vortex center. Note that $\beta \equiv (l_a)/l_\infty$, where l_a is the air parcel's angular momentum at the radius of maximum wind speed. Alternatively, we can write

$$d_a = \frac{\beta d_\infty \left(v_\infty + \frac{1}{2} f_\infty d_\infty \right)}{v_a}. \tag{11}$$

Since d_∞ is the initial distance from the vortex center of air parcels spiraling in toward the center, we refer to it as the *radius of influence* of the vortex. For dust devils and waterspouts, the earth's angular velocity can also be neglected at large radius. It follows from Eqs. (8) and (10) that the radius of maximum wind of these convective vortices is

$$a \approx \frac{\beta l_\infty}{\sqrt{RT_\infty \left\{ 1 - \exp \left[\left(\frac{\gamma\eta}{\gamma\eta - 1} \right) \left[\left(\frac{c_p}{R} \right) \left(\frac{T_0 - T_\infty}{T_\infty} \right) + \left(\frac{L_v}{R} \right) \left(\frac{r_0 - r_\infty}{T_\infty} \right) \right] \right\}}}. \tag{12}$$

Equation (12) shows that the vortex's radius of maximum wind depends only on thermodynamic variables, the air parcels' angular momentum at the vortex's radius of influence, and the fractional loss of angular momentum. Naturally, in order for the tangential momentum of air parcels spiraling in from opposite sides of the

vortex not to cancel each other, the ambient wind field must have *nonzero* vorticity. Therefore, in order to apply Eq. (12) to an ensemble of air parcels spiraling in toward the vortex center, l_∞ should be interpreted as the circulation (divided by 2π) around the vortex's radius of influence. Thus, Eq. (12) can be rewritten as

$$a \approx \frac{\frac{\beta}{2\pi} \oint \mathbf{v} \cdot d\mathbf{l}}{\sqrt{RT_\infty \left\{ 1 - \exp \left[\left(\frac{\gamma\eta}{\gamma\eta - 1} \right) \left[\left(\frac{c_p}{R} \right) \left(\frac{T_0 - T_\infty}{T_\infty} \right) + \left(\frac{L_v}{R} \right) \left(\frac{r_0 - r_\infty}{T_\infty} \right) \right] \right\}}}, \tag{13}$$

where \mathbf{v} is the vector velocity (absolute), and $d\mathbf{l}$ is the incremental distance along the integral path. The integral $\oint \mathbf{v} \cdot d\mathbf{l}$ is performed around a closed curve lying at the vortex's radius of influence (d_∞). That is, the circulation is computed around the vortex's radius of influence. It follows from Stoke's theorem that this cir-

ulation is equal to the integral of the vorticity over the vortex (the surface bounded by the curve at d_∞). Therefore, Eq. (13) shows that the vortex's radius of maximum wind depends only on the thermodynamics of the convective heat engine and the value of the background vorticity.

It follows from Eq. (13) that increases in the ambient vorticity lead to increases in the vortex's radius of maximum wind. This statement is valid for convective vortices in general. Torques due to friction lead to decreases in the angular momentum of air parcels spiraling in toward the vortex's center ($0 \leq \beta \leq 1$). Thus, with $\beta \approx 1$, Eq. (13) provides an upper bound for the radius of maximum wind. Since friction is larger near the surface, we should expect the vortex radius to be smaller near the surface. Indeed, this is what is observed in convective vortices.

3. Observations and discussion

To a first approximation, both the updraft and the downdrafts associated with waterspouts are saturated. From this approximation and from Golden (1974a), Simpson et al. (1986), and Golden and Bluestein (1994) observations, it follows that for a weak waterspout $T_\infty \approx 299.5$ K, so that $r_\infty \approx 0.022$; $T_0 \approx 300.5$ K, so that $r_0 \approx 0.024$; and the top of the convective layer is at $z_{\text{top}} \approx 3$ km, which corresponds to $\eta \approx 0.1$. Moreover, Rennó et al. (1998) showed that $\gamma \approx 1$. Therefore, we get from Eqs. (3) and (8) that for a weak waterspout $\Delta p \approx 6.5$ hPa, and $v_a \approx 25$ m s⁻¹. Since, by definition, a waterspout must have a spray ring, and a spray ring forms only when the wind speed is above $v_a \approx 20$ m s⁻¹, our result is consistent with the observations. For a strong waterspout, we have $T_\infty \approx 298$ K, so that $r_\infty \approx 0.020$; $T_0 \approx 302$ K, so that $r_0 \approx 0.025$; and $z_{\text{top}} \approx 10$ km, which corresponds to $\eta \approx 0.2$. It follows from equations (3) and (8) that $\Delta p \approx 40$ hPa, and $v_a \approx 60$ m s⁻¹. This result is also consistent with the observations.

Now, we use the heat engine framework to interpret some of the qualitative features of waterspouts. Golden (1974a) identified the dark spot as the initial stage of a waterspout formation. Initially, air parcels moving toward the convective plume absorb sensible and latent heat from the ocean while the vortex intensifies. During the dark spot stage, the vortex tangential wind speed reaches ~ 10 m s⁻¹. The plume's intensification, in turn, produces an increase in the surface heat flux. This positive feedback continues unabated until the vortex reaches the maximum potential intensity predicted by our theory or until it is disrupted by other means. Golden (1974a,b) observed that the majority of waterspouts do not evolve past this first stage. Golden suggests that this occurs because the incipient vortices are disrupted by either cool outflows from nearby showers or downdrafts from the parent cloud.

It has been suggested by Golden that the source of vorticity for the waterspout vortex is the wind shear between the warm updraft inflow and the rain-cooled downdraft outflow. That is, the source of vorticity is solenoidal. This idea is supported by the observations that waterspouts generally form over warm waters, close to their boundary with colder waters (Golden 1974b;

Peterson 1978; Simpson et al. 1986). However, frequently a rain-cooled downdraft outflow cuts off the source of warm air, and the waterspout decays.

Golden (1974a) observes that most waterspouts do not intensify past the spray ring stage, which corresponds to a tangential wind speed of about 20 m s⁻¹. These vortices are classified by Golden as weak waterspouts. We propose that most waterspouts do not intensify past the weak stage because the spray evaporation provides a negative feedback by cooling the updraft air. In cases where the convective heat engine is potentially powerful, however, the condensation funnel may reach down the surface shortly after the spray ring forms. The saturation of the near-surface air then nullifies the negative feedback by suppressing the spray evaporation. These waterspouts are then capable of intensifying until they either reach their full potential intensity or until they are destroyed by some other mechanism.

The idea that most waterspouts do not intensify past the weak stage because of the spray evaporation is supported by the following observations. (i) The evaporation of water can potentially cool the air down to its wet-bulb temperature; (ii) the convective updrafts are at most a few kelvins warmer than the ambient air; (iii) over the tropical oceans the wet-bulb temperature of the near-surface air is, in general, a few kelvins lower than its temperature.

4. Conclusions

We showed that the scaling theory for dust devils proposed by Rennó et al. (1998) is also successful in predicting the intensity of waterspouts. Moreover, we showed that it provides simple physical explanations for many of the general characteristics of waterspouts. It follows from our theory that the potential intensity of convective vortices is a function solely of the thermodynamic properties of their environment. Thus, given the environmental conditions, the potential intensity of convective vortices can be calculated.

The heat engine framework can be used to estimate the regions where these convective vortices are more likely to form. Equation (8) shows that the vortex intensity depends on the difference in temperature and water vapor content of the air at large radius (∞) and at the center (\odot) of the convective vortex. This, in turn, can be interpreted as the difference in temperature and water vapor content between updrafts and downdrafts (Rennó and Ingersoll 1996). In fact, this difference determines the heat input into the convective heat engine. Thus, convective vortices are more likely to form in the regions where the occurrence of the warmest and most moist updrafts and the coldest and driest downdrafts are supported by the local environment. These are the regions where the heat input into the convective heat engine is maximum and the solenoidal generation of vorticity is the greatest. This explains why waterspouts are frequently observed near the boundaries between rela-

tively warm and relatively cold waters (Golden 1974a). Since the work done by the convective heat engine is equal to the total heat input multiplied by its thermodynamic efficiency, another necessary condition for the formation of a most intense vortex is the presence of deep convection or developing deep convection.

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