Initiation of Run-Out Flows on Venus by Oblique Impacts

Seiji Sugita
Department of Earth and Planetary Science, Graduate School of Science, University of Tokyo, Hongo, Bunkyo-ku, Tokyo 113-0033, Japan
E-mail: sugita@eps.s.u-tokyo.ac.jp

and

Peter H. Schultz
Department of Geological Sciences, Brown University, Providence, Rhode Island 02912

Received March 31, 2000; revised August 16, 2001

A hypervelocity oblique impact results in a downrange-moving vapor cloud, a significant fraction of which is derived from the projectile. Since the vapor cloud expands to great extent and becomes very tenuous quickly on a planet with a thin or no atmosphere, it does not leave a well-defined geologic expression. The thick atmosphere of Venus, however, is sufficient to contain such a rapidly expanding vapor cloud. As a result of atmospheric interactions, impact vapor condenses and contributes to run-out flows around craters on Venus.

Previous results of both laboratory experiments and simple semi-analytical calculations indicate that an impact-vapor origin can account for the morphology of run-out flows on Venus most consistently. However, the detailed dynamics and geologic record of downrange-moving impact vapor clouds in Venus's atmosphere are not understood quantitatively. To approach these problems, we carried out two-dimensional hydrocode calculations. Parametric studies of these hydrocode calculations yield simple scaling laws for both the total downrange travel distance and the final temperature of impact vapor clouds under conditions on Venus.

Under typical impact conditions, impact vapor clouds travel downrange more than a crater radius prior to the completion of crater formation. Furthermore, the scaling law for the total travel distance is compared with observations for the downrange offset of the source regions of run-out flows around oblique craters. The results of this comparison suggest that energy/momentum-partitioning processes other than pure shock coupling may play important roles in hypervelocity impact at planetary scales. The results of hydrocode calculations also indicate that the terminal temperature of the impact vapor is close to the condensation temperatures of silicates, suggesting that two scenarios are possible for expected range of impact conditions:

1. Impact vapor condenses and forms run-out flows.
2. Impact vapor fails to condense and leaves no run-out flows.

Consequently, natural variation in impact angle, velocity, and projectile composition may account for partial occurrence of run-out flows around impact craters on Venus. © 2002 Elsevier Science (USA)

Key Words: cratering; impact processes; Venus; surface.

1. INTRODUCTION

1.1. Impact Vapor Clouds

Laboratory experiments (e.g., Lange and Ahrens 1982, 1986, Tyburczy and Ahrens 1986, Schultz 1988, Schultz and Gault 1990, Schultz 1996), hydrocode calculations (e.g., O’Keefe and Ahrens 1977, 1982, Orphal et al. 1980, Pierrazzo and Melosh 2000), and scaling relations due to dimensional analysis (e.g., Dienes and Walsh 1970, Schmidt and Housen 1987) predict that a large-scale hypervelocity impact generates a significant amount of impact vapor. A large portion of a silicate impactor is expected to evaporate at typical planetary-scale impact velocities 25–35 km/s (e.g., O’Keefe and Ahrens 1977). On planets with a thin atmosphere, such a large mass of impact vapor is expected to either escape from the atmosphere or expand to such a degree that it does not leave a well-defined and easily distinguishable signature in the geologic record (Melosh and Vickery 1989, Vickery and Melosh 1990).

The very thick atmosphere on Venus (90 times the mass of that of the Earth), however, may contain and prevent the impact vapor cloud from escaping (Schultz 1992c). A semi-analytical approach by Vickery and Melosh (1990) shows that almost all of the vaporized impactor material should be retained even on the Earth for impactor diameters and impact velocities smaller than 10 km and 20 km/s, respectively. Consequently, the vaporized impactor will be trapped in Venus's atmosphere and leave significant geologic signatures on the surface of Venus. Schultz (1992c) proposed that the run-out flows around impact craters on Venus revealed in Magellan images are related to the atmospherically decelerated and confined impact vapor and melt and showed that the morphologies of the geologic features are consistent with laboratory experiments and first-order calculations.

1 On leave at NASA Ames Research Center, MS245-3, Moffett Field, CA 94035. Fax: (650) 604-6779. E-mail: ssugita@mail.arc.nasa.gov.
20N334). The diameter and location of the crater are 31 km and 21.5° N, 332° E, respectively. Radar-bright run-out flows are seen on the lower right side of the crater.

1.2. Run-Out Flows on Venus

Precise radar mapping by the Magellan spacecraft revealed that impact craters on Venus are often associated with run-out flows such as that shown in Fig. 1 (e.g., Phillips et al. 1991). Although proposed origins of run-out flows on Venus differ, interpretations of morphological features of the run-out flows are generally consistent among authors (Phillips et al. 1991, Schultz 1991, 1992c, Schaber et al. 1992, Asimow and Wood 1992, Chadwick and Schaber 1993). Run-out flows are unique to Venus and are not found on any other terrestrial planets or satellites (Phillips et al. 1991). Although ejecta/melt flows on crater rims do occur on other planets (Schultz 1976, Hawke and Head 1977), their morphologies are very different from that of run-out flows on Venus (Schultz 1992c). Thus their formation process must be unique to Venus. Nevertheless, run-out flows do not occur around all craters on Venus. They tend to occur around craters with asymmetric morphology (i.e., due to oblique impacts) and larger diameter (Schultz 1992c, Schaber et al. 1992, Chadwick and Schaber 1993). The occasional occurrence of run-out flows is difficult to explain simply by impact melt, because large amounts of impact melt should be produced in any large-scale impact at velocity higher than the escape velocity of Venus (e.g., O’Keefe and Ahrens 1982). Yet less than a half of the craters smaller than 30 km in diameter are associated with run-out flows (Schaber et al. 1992). The role of impact vapor could, however, readily explain this observation. We argue below that impact vapor from some types of impacts on Venus may not reach the condensation condition before buoyancy carries it to the upper atmosphere. These cases will not produce run-out flows on the surface of Venus.

More subtle morphologies of the run-out flows also provide crucial information about their origin. The initial direction of run-out flows near their source regions appears to be more consistent with impact direction rather than local topography, which appears to control flow direction of the distal portion of run-out flows (Schultz 1991, 1992c, Asimow and Wood 1992, Chadwick and Schaber 1993). Stratigraphic relations show that emplacement of run-out flows occurs prior to that of lobate ejecta, although this stratigraphic relation is not clear around some craters (Schultz 1991, 1992c, Phillips et al. 1991, Asimow and Wood 1992). These observations are not consistent with a theory that clastic ejecta was fluidized and served as the source of run-out flows. Clastic ejecta is generated in the late excavation stage and deposits radially symmetrically, not concentrated along the trajectory direction as early-stage products, such as impact melt and vapor (e.g., Gault and Wedekind 1978, Schultz and Gault 1990). The deposits of clastic ejecta on a planet with an atmosphere are also limited in extent due to vortex entrainment and deceleration (Schultz 1992b, Barnouin-Jha and Schultz 1996), in contrast with the long run-out flows on Venus.

The near-source portion of many run-out flows have very diffuse boundaries (Schultz 1992c). When a run-out flow collides with a topographic obstacle, it sometimes splits into two flows: a flow that is stopped by the relief and a flow that overrides the relief. The blocked flow usually has lobate boundaries; the other has diffuse ones (Schultz 1992c). This variety in morphological appearances of run-out flows is also difficult to explain with a simple impact melt origin, but it can be explained readily by suspension flows or debris flows analogous to pyroclastic flows, which consists of hot impact vapor, impact melt, impact vapor condensate, and entrained ejecta (Schultz 1992c). As a suspension flow advances, it deposits liquid/solid phases. If the deposited material is dominated by the liquid phase, it can coalesce to form a liquid, thereby producing a lava-like flow (Schultz 1992c). This transition also may explain rather equivocal results of rheological analyses of run-out flows that lead to estimates of the viscosity of run-out flow materials to have a wide range of values: from $<10^{-3}$ Pa s (~suspension flows) to $10^5$ Pa s (~basaltic lava flows) (Asimow and Wood 1992, Johnson and Gaddis 1996). In fact, the high atmospheric pressure of Venus allows condensation of impact vapor at high temperature, which makes it more likely for a liquid phase to be dominant rather than a solid phase. It also should be noted, however, that some of the low viscosity values may not be necessary to account for the flow patterns if the mass discharge rate of the flows are very large (Miyamoto and Sasaki 2000).

Among various possible origins for run-out flows, only the impact-vapor-cloud scenario appears to be able to account for all the above observations. Simplified analytical models by Schultz (1992c) also indicate that both downrange deceleration and
atmospheric containment of impact vapor are significant on Venus and generally consistent with geologic observations. The vapor cloud hypothesis, however, still raises crucial questions regarding consistency with geologic observations. 1. Will the impact vapor cloud reach the final crater rim prior to the ejecta curtain? 2. Will impact vapor actually condense before ejecta emplacement? If not, then a run-out flow cannot be initiated prior to ejecta. To answer these questions, we need to carry out hydrodynamic calculations to describe the time-evolving conditions of the vapor controlled by shock interaction with the ambient atmosphere.

We then proceed one further step and ask: What is the geologic significance of run-out flows on Venus? In this study, we explore the significance of both downrange offsets of the source regions of run-out flows and occasional occurrence of run-out flows around impact craters.

The calculations are carried out for a variety of combinations of controlling parameters such as initial conditions of a vapor cloud. Such parametric studies reveal which variable controls which aspect of a run-out flow. Finally, the results of the hydrodynamic calculations are compared with a quantitative measurement of run-out flows on Venus in order to understand their physical significance.

2. NUMERICAL EXPERIMENTS

2.1. Modeled Impact Vapor Clouds

Laboratory experiments indicate that downrange-moving vapor clouds induced by oblique impacts (generally less than about 30° impact angles measured from the horizontal) travel along the ground surface (Schultz 1996). Consequently, we model the initial condition of a downrange-moving vapor cloud as a gas body with translational momentum along the ground surface. The geometric configuration of a modeled impact vapor cloud is illustrated in Fig. 2. To focus on the main questions raised in the previous section, we assess only atmospheric deceleration of translational motion and radial expansion of downrange-moving vapor clouds. Other aspects of the dynamics of vapor clouds such as buoyancy rise, gravitational collapse, and basal drag on the ground surface are omitted here. Such processes affect only the last stages of emplacement. We also assume for calculational convenience that the shape of the impact vapor cloud is a hemisphere initially. Then the calculation is reduced to the dynamics of an initially hemispherical gas body with both high temperature/pressure and translational motion within a homogeneous atmosphere. The numerical experiments in this study are unique in that they focus on the horizontal motion of an impact vapor cloud within an atmosphere, whereas previous studies have looked at either vacuum cases (e.g., O’Keefe and Ahrens 1986) or vapor clouds due to vertical impacts (e.g., Vickery and Melosh 1990).

Temperature, density, and translational velocity are assumed to be initially homogeneous within the gas body. The initial radial expansion velocity is assumed to be zero. Because this dynamic system is cylindrically symmetric, the hydrodynamic equations to describe the dynamics are written as

momentum equation (Euler equation)
\[
\frac{du}{dt} = -\frac{1}{\rho} \nabla p, \tag{1}
\]
continuity equation (compressible flow)
\[
\frac{dp}{dt} = -\rho \nabla \cdot u, \tag{2}
\]
energy equation
\[
\frac{de}{dt} = -\frac{p}{\rho} \nabla \cdot u, \tag{3}
\]
equation of state (ideal gas)
\[ p = (\gamma - 1) \rho \frac{e}{\gamma} \]

where \( t, u, \rho, p, e, \) and \( \gamma \) are time, velocity vector, density, pressure, specific internal energy, and the ratio of specific heats, respectively. The other notations are defined as

\[ \frac{d}{dt} \equiv \frac{\partial}{\partial t} + u \frac{\partial}{\partial x} + v \frac{\partial}{\partial r}. \]

\[ \nabla \cdot u = \frac{\partial u}{\partial x} + \frac{1}{r} \frac{\partial r}{\partial r}, \]

\[ \nabla p = \left( \frac{\partial p}{\partial x} \frac{\partial p}{\partial r} \right). \]

where \( x, r, u, \) and \( v \) are the coordinate in the direction of the translational motion, radius from the \( x \)-axis, velocity in the \( x \)-direction, and radial velocity, respectively. The origin of the \( x-r \) coordinates is taken at the initial mass center of the gas body (Fig. 2). For intuitive interpretation of calculation results, the origin of the coordinate system and the \( x \)-axis may be considered as the impact point and the direction of downrange motion of the resulting vapor cloud along the ground surface, respectively. Such a geometric configuration is the closest approximation to a downrange-moving vapor cloud due to an oblique impact observed in laboratory experiments (Schultz 1992c, 1996).

The above set of equations describes shock phenomena in a compressible fluid and is often used in the literature (e.g., Richtmyer and Morton 1967, Woodward and Collela 1984, LeVeque 1992). It omits, however, three important processes: viscosity, radiation, and the non ideal nature of the gas. First, omission of a viscous term is justified because the Reynolds number \( (Re \equiv UD/\nu) \) of an impact vapor cloud is extremely large \( (>10^{14}) \) owing to the small kinetic viscosity of the heated gasses \( (\nu \sim 3 \times 10^{-7} \text{ m}^2/\text{s}) \) and the large diameter of the impact vapor cloud \( (D > 1 \text{ km}) \) and its large velocity \( (U \sim 10 \text{ km/s}) \).

Second, simple analytical assessments reveal that radiative cooling effects are negligible in comparison to adiabatic cooling in the early stage of vapor cloud evolution under the Venus condition (see Appendix). Third, the discrepancy between a real and ideal gas is small until the temperature and pressure approach the condensation condition. Our primary interest here is the earlier stage of evolution of the vapor cloud, not the condensation process itself. Consequently, the detailed dynamics of the condensing impact vapor clouds in the later stages is not explored in the present study. Another uncertainty in the equation of state (Eq. (4)) is \( \gamma \). The ratio \( \gamma \) of specific heats is large (e.g., 1.4 for linear molecules) at lower temperatures and is small (e.g., 1.1–1.2) at high temperatures. However, precise evaluation of \( \gamma \) increases the computational time drastically. This will prevent us from conducting a comprehensive parametric study for many controlling factors. A field observation of the extremely intense shock wave due to a nuclear explosion indicates that 1.2 provided a good approximation for \( \gamma \) for a wide range of temperature and pressure (Taylor 1951). Thus we use 1.2 for \( \gamma \) in this study. The air and impact vapor are assumed to be \( \text{CO}_2 \) \( (\mu = 44 \text{ g/mol}) \) and forsterite vapor \( (2\text{MgO} + \text{SiO}_2 + \frac{1}{2}\text{O}_2, \mu = 40 \text{ g/mol}) \), where \( \mu \) is molecular weight.

There also may be concerns for the effects of the ground surface and nonspherical shape of an initial vapor cloud. If all the aspects of evolution of a cloud were to be reproduced, such effects would have to be considered. It is often confusing, however, to combine all the processes together before understanding each of them separately. Consequently, we confine ourselves here in a framework of testable numerical experiments and try to understand the basic shock phenomena produced by an idealized moving/expanding gas body within an atmosphere.

2.2. Numerical Procedure

A two-dimensional hydrocode was written based on the CIP (Cubic Interpolated Propagation) method of Yabe and Aoki (1991) and Yabe et al. (1991). CIP is a numerical scheme to
FIG. 3—Continued
calculate advection and to capture shock fronts with third-order accuracy in space. Advantages of the CIP method are its simplicity and computational efficiency compared to other algorithms with the same order of accuracy (see, e.g., Woodward and Colella 1984, LeVeque 1991). This allows one to carry out many of the large runs on personal computers. CIP also has many other excellent advantages as discussed elsewhere (Yabe 1992).

Hydrocodes based on the CIP method have been successfully applied to Shoemaker-Levy-9 impacts into Jupiter as well as other fields of physics, thereby demonstrating its effectiveness (e.g., Yabe 1992, Yabe et al. 1994, 1995).

Because the CIP algorithm calculates only the advection part of the hydrodynamic equations, an additional algorithm is needed to compute nonadvective terms. Our code adopts the finite difference scheme of Yabe et al. (1991) for these nonadvective terms. The code differs from that of Yabe et al. (1991), however, in the time-integration method. Our code uses the so-called backward-Euler iteration (e.g., Kurihara 1965) to increase the stability of the code, which allows the code to calculate stably an extremely large pressure contrast (> 10^4 times) across the shock wave around an expanding impact vapor cloud.

Because the size of a vapor cloud changes by more than an order of magnitude from its formation to the final stage (Schultz 1992c, Sugita and Schultz 1995), it would be extremely inefficient (and inaccurate) to use the same uniform mesh size throughout a run. Thus our code was designed to rezone the mesh when the outermost shock wave reaches an edge of the calculation domain. The re zoning is done in the x-direction and the r-direction individually, upon arrival of a shock wave at a respective outer edge of the calculation domain. This two-step approach reduces changes in the size of the calculation grid due to a re zoning. Each re zoning doubles the length of the calculation grid in one direction and keeps the length in the other direction the same. Since the total number of grid points is kept constant, the new calculation domain covers twice the space. An impact vapor cloud produced by an oblique impact also has a translational velocity downrange (along the x-axis). At the end of each time step, the locations of downrange-moving and uprange-moving shock fronts are identified, and the mesh points are shifted such that the geometric center of the shock front circle is always located at the center of the calculation domain. This procedure maximizes efficiency while it reduces the number of rezonings in the x-direction.

Most of the calculations here were done on a numerical grid space with a size of 256 × 128, where the larger dimension is taken along the x-axis. The reflective boundary condition is imposed along the axis of symmetry (i.e., at \( r = 0 \)). To assess the effect of the size of the calculation grid, we carried out calculations on grid sizes of both 512 × 256 and 128 × 64. Although specific hydrodynamic features are affected by the change in the grid size, the general trends and bulk motion of calculated vapor clouds are not influenced.

It is emphasized that the calculated system is nondimensionalized with the initial size of a vapor cloud. But for convenience of intuitive interpretation of the results, all the scales in the figures that follow are labeled with absolute scales for a nominal case, in which the initial radius of a vapor cloud is 1 km.

3. CALCULATION RESULTS

3.1. Hydrodynamic Phenomena

The extremely large pressure gradient across the material boundary between an impact vapor cloud and the ambient atmosphere induces a very strong outward radial flow (Fig. 3c). As this flow collides with the ambient atmosphere, it creates a very intense shock front around the vapor body. The flow velocity colliding with the ambient air is much greater in front of the cloud than behind because of its rapid downrange translational motion. This leads to an asymmetric shock front that is stronger in the front and weaker in the back (Figs. 3a and 3b). The asymmetric shock wave can be viewed as an intermediate form between a bow shock and point-source explosion.

The evolving shape of the shock wave is a function of both the flow speed of expanding vapor and the translational velocity of the vapor cloud. When the expanding flow is relatively strong, the shape of the shock front becomes blunt. But when the translational velocity is very high, the shock front becomes more streamlined. Thus the cross-sectional surface area of a vapor cloud with high translational velocity becomes smaller, thereby reducing the effect of air drag on cloud deceleration. This phenomenon resembles the shock collimation of debris material passing through a thick atmosphere (Schultz 1992c, Schultz and Sugita 1994).

The outgoing flow carries vaporized material away from the high-density core of the cloud and pushes it against the leading hemispherical shock front. In Fig. 3d, the high-density vapor core has disappeared. Most of the vapor mass is in the hemispherical shell just behind the shock front.

Figures 3d and 3e reveal small ripples on the interface between the impact vapor and the ambient atmosphere. This rippling is a Rayleigh–Taylor (R–T) instability as often seen in a variety of impact phenomena (e.g., O’Keefe et al. 1994). The consequence of R–T instabilities on the evolution of the impact vapor cloud is discussed in the next section. The instability, however, does not evolve too far in the calculations because it is constantly stretched and dragged away by ambient flow around the air/vapor boundary. Nevertheless, the detailed dynamics of this boundary surface may be more complex than shown in the macroscopic calculations because it may be controlled by fine-scale processes such as local eddies much smaller than the calculation resolution. The detailed interaction between such microscopic R–T instabilities and macroscopic phenomena, however, is beyond the scope of this study and will be explored as a separate contribution.

After the high-density core of vapor disappears, a backward-moving secondary shock front forms and detaches from the first shock front (Figs. 3g, 3h, and 3i). This secondary shock front propagates through the vapor cloud and then creates a plume directed backward from the cloud. As this plume grows farther,
it induces a Kelvin–Helmholtz (K–H) instability that leads to a mushroom-shaped cloud (Figs. 3j, 3k, 3m, and 3n). The mushroom cloud occurred in almost all the calculated cases. Thus, it appears to be a characteristic product of the process that may occur in more complex real three-dimensional vapor clouds.

The pressure in the vapor cloud eventually reaches equilibrium with the ambient atmosphere, and the vapor cloud ceases expansion. The shock wave generated by the vapor cloud, however, continues its expansion and leaves the vapor cloud behind (Figs. 3n and 3o). When this shock detachment occurs, the vapor cloud typically has grown about 10 times its initial radius while the downrange translational velocity has reduced to 1/10 of its initial value or even less. Because strong shock waves inside the vapor cloud have disappeared at this stage, the effects of the subsonic flow field become important. Figure 3o shows the flow field around a late-stage impact vapor cloud and reveals development of strong vortices that last well after the shock waves have dissipated.

The distribution of temperature in the vapor cloud is very heterogeneous in the early stage: It is very high in the shock front but relatively low in the region behind the shock (Figs. 3b and 3e). Also note that compressed ambient air has a much higher temperature than that of the inner impact vapor. After shock detachment, the extremely hot compressed air is left behind the impact vapor cloud, and the relatively cooler impact vapor moves farther downrange (Figs. 3k and 3n).

### 3.2. Time Evolution

Both the center of mass and the average temperature of each vapor cloud are monitored during the calculations. The motion of the center of mass reflects cloud deceleration due to air drag. The average temperature measures the rate of the adiabatic cooling and secondary heating due to shock waves created by collision between a vapor cloud and the surrounding atmosphere. The combined data provide a basis for assessing possible observable products recorded on the surface of Venus.

#### 3.2.1. Mass center motion

Although the motion of the mass center of a vapor cloud is quite complex, its overall pattern is always the same and does not depend on specific values of parameters for initial condition, such as specific internal energy $\epsilon_{\text{vap}}$ and translational velocity $V_{tr}$ (Fig. 4). A simplified air-drag model is useful for understanding the general behavior of the center-of-mass motion,

$$M \frac{dU}{dt} = -\frac{\pi}{2} C_D \rho_{\text{air}} R^2 U^2,$$

where $M$, $R$, $U$, $t$, $C_D$, and $\rho_{\text{air}}$ are the mass and radius of a vapor cloud, the velocity of the mass center, time, drag coefficient ($\sim0.5$), and air density, respectively. This equation holds well for an object going through a uniform atmosphere when the Reynolds number is large (e.g., Landau and Lifshitz 1987). When the size $R$ of the vapor cloud is kept constant, Eq. (8) can be solved analytically. Schultz (1992c) used this classical solution to assess the deceleration process of impact vapor clouds in Venus’s atmosphere, showing that a vapor cloud needs to have a large fraction of impactor mass to travel beyond the crater rim. Figure 4 shows that the constant-size solution approximates well a hydrocode calculation during the early stages. The constant-size model, however, cannot adequately describe the motion of the vapor cloud at later stages, when it predicts too low a deceleration rate for a vapor cloud. This problem results from the expansion of a vapor cloud. As its cross-sectional surface area ($\sim R^2$) increases, the drag force rapidly increases and stops the cloud motion more efficiently. Consequently, a 1-km-radius cloud travels 10 km within the first several seconds and virtually stops before 10 s after its formation. Scaling relations based on laboratory experiments (Schmidt and Housen 1987) predict that a crater formed by such an impact should take about 30 s to complete the excavation stage. Ejecta emplacement is expected to be further delayed because of large atmospheric drag and vortex formation by the moving ejecta curtain (Schultz and Gault 1982, Schultz 1992b, 1992c, Barnouin-Jha and Schultz 1996). The downrange motion of the vapor cloud created by a kilometer-size projectile, therefore, will be completed well before ejecta emplacement.

After traveling downrange to its maximum downrange distance, however, the vapor cloud does not simply stop. When
its translational velocity reaches zero, it actually starts going backward (i.e., uprange). Then it stops again and starts moving downrange. Such motion repeats a number of times, thereby developing oscillation (Fig. 4). This oscillatory motion is due to both vortex motion around a vapor cloud and shock/acoustic vibration within it. Although the amplitude of this oscillation is much less than the first downrange movement, it occurs over a comparable time and is the last process interacting with the surface.

Figure 4 also compares the maximum buoyancy rise/fall rate of the impact vapor cloud with the result from the hydrocode calculations. The maximum rise/fall is simply given by the gravitational free fall formula, \( \frac{1}{2}gr^2 \), where \( g \) and \( t \) are gravity and time, respectively. The comparison clearly shows that the buoyancy effect is very small in the early stage of evolution but becomes extremely important in the late stage.

### 3.2.2. Temperature history

The mean temperature is defined as the mass-weighted average of temperature of all the impact vapor in a cloud. Because the temperature structure is not homogeneous in such a cloud (as shown in Fig. 3), some portions will be significantly hotter or cooler than the mean temperature. It is particularly important to note that condensation may start in certain regions well before the mean temperature goes below the condensation temperature.

Condensation, however, is problematic for characterizing the temperature in this study. Because the ideal-gas equation of state is assumed, the thermodynamic relations among pressure, density, temperature, and others are not accurately calculated around conditions of condensation. Moreover, the dynamics of the cloud will be further affected by significant heat release during the condensation process. Thus the detailed dynamics around the condensation condition may not be adequately predicted by the model. It is also difficult to judge if condensation actually occurs due to adiabatic cooling. The general dynamics well before vapor condensation, however, should not suffer from this simplification. Similarly, if the calculated temperature goes well below the condensation temperature of the impact vapor, we can safely assume that condensation will actually occur despite the ambiguity.

Figure 5 shows that the mean temperature of a vapor cloud initially decreases rapidly to a terminal temperature. The time scale for this temperature drop is approximately the same as the deceleration of the downrange motion. In some calculations, the terminal temperature is well below the probable condensation temperature of silicate rocks and metals in a 90-bar condition (3000–4000 K; see Thompson and Lauson (1972), Stevenson (1987), and Benz et al. (1989)). In these specific cases, most of the impact vapor in the cloud will condense well before crater formation and ejecta emplacement. This is consistent with the observed stratigraphic relations indicating that the emplacement of many run-out flows occurs prior to that of ejecta (Schultz 1991, 1992c, Phillips et al. 1991, Asimow and Wood 1992).

### 3.3. Parametric Study

A series of parametric studies were performed to understand how the total travel distance and the terminal temperature of impact vapor clouds depend on both the initial conditions of the vapor clouds and the ambient atmosphere. The parameters used were the initial specific internal energy \( \varepsilon_{\text{vap}} \), the downrange translational velocity \( V_{\text{tr}} \), density \( \rho_{\text{vap}} \), and mass \( M_{\text{vap}} \) of a vapor cloud and the ambient atmospheric density \( \rho_{\text{air}} \). A number of calculations were done for all these conditions. Such analyses reveal what conditions enable an impact vapor cloud to travel beyond the crater rim and to condense to form the observed run-out flows.

#### 3.3.1. Total travel distance

The total travel distance \( L \) here does not refer to the actual terminal position of a vapor cloud but to the first maximum downrange excursion in the oscillatory motion (see Fig. 4). There are two reasons for this choice. First, the resolution in the calculations becomes increasingly poor after shock detachment when the effective number of calculation cells to cover the vapor cloud becomes very small due to rezoning; the mesh size is 250 m near the end of the calculation. Thus the position of the mass center becomes very inaccurate during the very late stages in the calculations. The second reason involves the timing of condensation onset. Because the mean temperature of a vapor cloud reaches very close to its lowest value when the vapor cloud is at its first downrange maximum
position, condensation most likely occurs around this time. Thus the first maximum distance serves as a useful reference point for comparison with observed run-out flows on Venus.

Figure 6 reveals that normalized total travel distance $L$ of a vapor cloud can be expressed by a power-law function for both the initial specific internal energy $\varepsilon_{vap}$ and the initial downrange velocity $V_{tr}$ of a vapor cloud. A similar parametric study has been conducted for both ambient air density $\rho_{air}$ and the initial impact vapor density $\rho_{vap}$ over the range of $10 \leq \rho_{air} \leq 100 \text{ kg/m}^3$ and $1000 \leq \rho_{vap} \leq 8000 \text{ kg/m}^3$. The results indicate that $L/r_p$ is also approximated well by a power-law function of both $\rho_{air}$ and $\rho_{vap}$. The effects of all the variables are summarized by the semi-empirical scaling relation.

$$L/r_p = 13 \left( \frac{\rho_{air}}{67 \text{ kg/m}^3} \right)^{1/2} \left( \frac{\rho_{vap}}{3 \text{ g/cm}^3} \right)^{0.4} \left( \frac{M_{vap}}{M_{proj}} \right)^{1/3} \times \left( \frac{\varepsilon_{vap}}{50 \text{ MJ/kg}} \right)^{-1} \left( \frac{V_{tr}}{10 \text{ km/s}} \right)^{1.3}, \tag{9}$$

where $M_{vap}$ and $M_{proj}$ are the masses of vapor cloud and impactor, respectively. Here, the fourth factor $(M_{vap}/M_{proj})^{1/3}$ in Eq. (9) was introduced to normalize $L$ with projectile radius $r_p$ instead of the initial vapor cloud radius. This choice was made because $r_p$ can be estimated from geologic observation of Venus as discussed in Section 4.1. Note that the power-law index of the $V_{tr}$ term in Eq. (9) is larger than unity. This is because air drag is reduced at higher velocities. This rather counterintuitive phenomenon results from the shock collimation effect described above. The strong shock in front of the vapor cloud reduces the cross-sectional surface area more efficiently at higher velocities.

Equation (9) can be rewritten with another set of variables to express its significance more clearly. For given values of impactor size, air density, and initial vapor density, Eq. (9) predicts the same downrange travel distance for different conditions of vapor clouds. For example, a fast and hot cloud, a slow and cool cloud, and a fast cool and small cloud all travel the same distance downrange. In other words, the three parameters $M_{vap}$, $\varepsilon_{vap}$, and $V_{tr}$ have many combinations that will yield the same total travel distance. The three parameters are not, however, independent of each other. Because the total amount of momentum and energy available for an impact vapor cloud are limited, there should be an anticorrelation between vapor mass and the temperature and velocity of a vapor cloud. For a given size and velocity of impact, a larger mass of impact vapor cloud should have smaller velocity and temperature (or specific energy). To express this relation more explicitly, partition ratios $\phi_{\text{energy}}$ and $\phi_{\text{momentum}}$ are introduced,

$$\phi_{\text{energy}} = \frac{M_{vap} \varepsilon_{vap}}{\frac{1}{2} M_{proj} V_{im}^2}, \tag{10}$$
$$\phi_{\text{momentum}} = \frac{M_{vap} V_{tr}}{M_{proj} V_{im}}, \tag{11}$$

where $M_{proj}$ and $V_{im}$ are projectile mass and impact velocity, respectively. Since the velocities $V_{tr}$ and $V_{im}$ have different angles (measured from the horizontal), Eq. (11) deviates from the standard definition of momentum ratio. This definition, nevertheless, is useful for comparison with the results of impact.
Equation (12) reveals that $L/r_p$ is virtually independent of the ratio of vapor mass $M_{vap}$ to projectile mass $M_{proj}$ when the partitioning ratios are kept constant. Both $\rho_{air}$ and $V_{im}$ are independent of impact angle $\theta$. The initial vapor density $\rho_{vap}$, which is essentially the same as the density of a shock-compressed projectile, does not change with impact angle $\theta$ very much, because peak density under impact shock is a weak function of peak pressure under high shock pressure. For example, a linear shock-particle velocity equation of state for basalt (Kieffer and Simonds 1980) predicts about a 25% difference in compression between a 90° (vertical) impact and a 15° impact at 28 km/s of impact velocity. Furthermore, since the downrange travel distance is proportional to only the 0.4th power of the initial vapor density, 25% of the density difference results in less than 10% of the difference in travel distance $L$. Consequently, the travel distance $L$ of a downrange-moving impact vapor cloud as a function of impact angle mostly reflects the initial energy/momentum coupling during hypervelocity oblique impacts.

3.3.2. Mean temperature. Figure 7 shows the terminal temperature $T_f$ of a vapor cloud as a function of initial specific internal energy $\varepsilon_{vap}$ and initial translational velocity $V_{tr}$. Both higher specific internal energy and higher translational velocity result in a higher terminal temperature. A similar parametric study indicates that the terminal temperature of a vapor cloud depends little on the density of the ambient air, the initial density of impact vapor, and the ratio of specific heats. Thus $T_f$ is essentially a function of only $\varepsilon_{vap}$ and $V_{tr}$. The data in Fig. 7, however, do not follow a power law as in Fig. 6. When $T_f$ is plotted against effective specific energy $\varepsilon_{eff}$ of a vapor cloud,

$$
\varepsilon_{eff} = \varepsilon_{vap} + \frac{k}{2} V_{tr}^2,
$$

all the points from the parametric calculation collapse onto a straight line in a log–log plot (Fig. 8). Here, $k$ is a constant with a value of about 0.2. Since the slope of the fit line is approximately unity, the terminal temperature of a vapor cloud can be approximated by a linear relation

$$
T_f = 3500 \text{ K} \times \frac{\varepsilon_{eff}}{45 \text{ MJ/kg}}.
$$

Although the underlying reason for the specific value of 0.2 for the coupling coefficient $k$ in Eq. (13) is not yet fully understood, its physical significance is clear. The coupling coefficient is the ratio of the initial kinetic energy of a vapor cloud converted to its internal energy through complex shock interaction. The rest of the initial kinetic energy of the cloud, i.e., 80%, is directly converted to the internal energy of the ambient atmosphere. Not all of the 20% of the kinetic energy converted to the internal energy of the vapor cloud, however, is maintained in the cloud throughout its dynamic history. Part of it is eventually transferred to the ambient atmosphere as work done against the atmosphere through the adiabatic expansion process together with the intrinsic internal energy of the vapor cloud (i.e., “indirect” energy transfer mechanisms from the impactor to the atmosphere (Schultz 1992c)).

Here we choose the condensation temperature of forsterite as a reference, although the condensation temperature depends

![Fig. 7](image7.png)  
**Fig. 7.** The terminal mean temperature of impact vapor clouds as a function of both the initial downrange translational velocity $V_{tr}$ and the initial specific energy density $\varepsilon_{vap}$ of impact vapor clouds, including vaporization energy (12 MJ/kg).

![Fig. 8](image8.png)  
**Fig. 8.** The terminal mean temperature of impact vapor clouds as a function of effective energy. The effective energy is defined as the sum of internal energy and 1/5 of kinetic energy of an impact vapor cloud. Note that the terminal temperature is well approximated by a linear function of effective energy.
significantly on the chemical composition of the vapor material. Because pressure in a vapor cloud is well equilibrated with the ambient atmosphere during late stages, the condensation temperature is calculated at the atmospheric pressure of Venus (i.e., 90 bars). A numerical extrapolation using ANEOS equations of state by Benz et al. (1989) indicates that forsterite at 90 bars condenses at around 3500 K. This corresponds to 45 MJ/kg for the initial effective specific energy of vapor (see Eq. (14) and Fig. 8). This specific energy is in the middle of our estimated range of effective specific energy of impact vapor clouds. Here, if the ratio of the impact energy partitioned to the internal energy of projectile is about 1/4 for vertical impacts (Gault and Heitowit 1963) and decreases as the 1.5th power of the sine of the impact angle for oblique impacts (Pierazzo and Melosh 2000), then the calculated range of specific internal energy (i.e., from 25 to 100 MJ/kg) corresponds to 23° to 90° impacts at an impact velocity of 28 km/s. Because there is significant uncertainty in estimates for the initial specific energy of an impact vapor cloud (e.g., Stevenson 1987, Schultz 1996, Sugita and Schultz 1999), it is difficult to predict when exactly condensation occurs. Nevertheless, the qualitative calculation result that vapor clouds due to oblique impacts on Venus may reach both a condensing condition and a noncondensing condition should not be affected by such uncertainty.

3.4. Validity and Limitation of the Calculations

It is important to clarify the nature of the calculations and to organize which aspects of the calculations are comparable to a real impact vapor cloud and which are not, because the calculations in this study are considerably simplified. First, it is emphasized that the nature of the calculations are not simulations of real impact vapor clouds but rather numerical experiments to understand selected key physical processes in complex vapor–atmosphere interactions. The calculation focuses on atmospheric deceleration of translational motion and radial expansion of impact vapor clouds. The calculated simple two-dimensional system (i.e., moving an initially homogeneous gas sphere within a uniform atmosphere) is to be viewed as a first step to understand these processes. Consequently, other factors that may be important in a real three-dimensional impact vapor cloud are omitted here. Those factors include atmospheric density gradient, gravity, buoyancy force, both geometry and heterogeneity of an initial vapor cloud, friction along the ground surface, and entrainment of ambient atmosphere into a vapor cloud.

The two-dimensional approximation, however, should not affect the general trends, such as deceleration rate and decrease in average temperature, and can be shown to have minimal effects on the calculation results. First, the effect of atmospheric density gradient is not very important in the early stage of evolution of impact vapor clouds of the size considered in this study. An impact vapor cloud typically has traversed 60–70% of its downrange travel by the time it has radially expanded by a factor of four. If the initial radius of an impact vapor cloud is 1 km, the density ratio in Venus’s atmosphere between the top and the center of the expanded vapor cloud is only 1.3. If the top and bottom portion of a vapor cloud were dynamically detached, this difference in atmospheric density would result in less than a 10% difference in total downrange travel distance \( L \) (Eq. (12)). Second, by the time either the buoyancy force or gravity becomes important, the downrange motion of an impact vapor cloud has essentially ceased, as shown in Fig. 2. Such factors then will not affect the total travel distance of the cloud. Third, as discussed in Section 3.1, most of the mass of an impact vapor cloud is initially concentrated into a thin shell-like structure behind the frontal bow shock (see Fig. 3d). This process greatly reduces the effect of initial geometry and heterogeneity of an impact vapor cloud.

Other omitted factors, however, might be significant. First, existence of the ground surface may deform the shape of the moving vapor cloud, thereby affecting the deceleration rate. Large differential velocities along the ground surface should generate intense vortices and enhance the effective viscosity at the base of the vapor cloud. This process may increase the deceleration rate of the cloud significantly as well. Second, if such intense vortices are formed around the vapor cloud, they will enhance mixing between impact vapor and ambient atmospheric gas. Such mixing would delay the condensation of impact vapor. However, to solve these problems of ground surface and vortex formation, different types of hydrocode calculations are required. Consequently, these problems are left as future research themes.

4. GEOLOGIC IMPLICATIONS

The results of the hydrocode calculations are now compared with Venus geology in order to understand their significance. First, a quantitative measurement of the downrange offset of source regions of run-out flows on Venus is compared with the calculated total travel distance of downrange-moving impact vapor clouds (i.e., Eq. (12)) and laboratory experiments. Then, the significance of the average terminal temperature of impact induced vapor clouds is discussed to assess the condensation conditions of impact vapor, as well as effects of physical processes not included in the calculations. We also discuss wind streaks observed around impact craters on Venus.

4.1. Downrange Offset of Source Regions for Run-Out Flows

If the emplacement of run-out flow material is controlled by the dynamic motion of an impact vapor cloud as inferred from morphology of run-out flows on Venus, then the downrange offset of the source regions for run-out flows should represent the distance that an impact vapor cloud could travel before its condensation/deposition. To compare this offset with numerical results, we measure this dimension as a function of impact angle.

The offset distance can be defined as the difference between the first point of impact and the terminal position of the mass center of a vapor cloud. The location of the point of impact is approximated by the uprange edge of the central structure of a crater, i.e., a central peak or a central ring (Figs. 1 and 9). This position is uprange from the geometric center of a final crater cavity.
and is also the intersection point for early-stage ejecta from an oblique impact on the Moon (Schultz and Anderson 1996). The terminal position of a decelerating vapor cloud is approximated by the run-out flow “source” region around a crater from which run-out flows appear to emerge. This inferred source region has several characteristic features (Schultz 1992c). First, the flow direction in a source region is consistent with impact trajectory, rather than the local topography. Second, they are radar bright. Third, their boundaries are usually very diffuse. And fourth, the flows typically appear to have interfered with downrange ejecta emplacement and precede ejecta emplacement elsewhere. We judge an area that satisfies at least three of the four conditions as a run-out-flow source region. A source region usually fits well within either a circular or an elliptical area (Fig. 9). We approximate the terminal position of the impact vapor cloud with the center of this ellipse.

The observed value of the offset distance is not actually convenient for comparison with numerical calculations owing to the effect of impactor size. With all other conditions the same, a larger impactor (i.e., greater mass) creates a larger vapor cloud and will travel a greater distance before air drag stops it. When the size of a vapor cloud is sufficiently smaller than the atmospheric scale height, the final travel distance of the vapor cloud is proportional to the radius of the initial vapor cloud (Eq. (12)), which is proportional to the impactor size. Consequently, impactor size needs to be determined. This can be estimated from two different observable geologic dimensions.

The first one is crater diameter. Scaling analyses allow crater diameter to be related to impactor size for a given impact velocity (e.g., Schmidt and Housen 1987) but require additional corrections for post-formation modification and possibly large atmospheric effects on crater growth (Schultz 1992a, 1992c). A second impactor-size-dependent dimension is the size of central structures of craters such as central peaks, pits, and peak rings (Schultz 1988, 1994). This method is based on the hypothesis that the formation of the central structures is controlled by the initial penetration process of the impactor into the target. This proposal has been examined with laboratory experiments (Schultz and Gault 1993) and with morphologic observation of craters on Mars (Schultz 1988), Mercury (Schultz 1988), Venus (Schultz 1992c), the Moon (Schultz and Anderson 1996), and Earth (Schultz 1994 Schultz and Anderson 1996). It accounts for the larger size of central structures for shallower impact angles and also helps in understanding other observations of their structure. The scaling relation adopted here, from Schultz (1994) is

\[
\frac{W}{2r_p} = 2.2 \left( \frac{V_{im}}{c_t} \right)^{2/3} \left( \frac{\rho_p}{\rho_t} \right)^{1/3},
\]

where \(W\), \(r_p\), \(V_{im}\), \(c_t\), \(\rho_p\), and \(\rho_t\) are the lateral dimension of a central structure, impactor radius, impact velocity, the sound velocity of target material, and the densities of projectile and target, respectively. The target sound velocity \(c_t\), impact velocity \(V_{im}\), and the density ratio \(\rho_p/\rho_t\) are assumed to be 5.5 km/s,
28 km/s, and 1, respectively. It is noted that Eq. (15) is applicable to craters created by oblique impacts but that it does not contain the impact angle $\theta$ explicitly. This advantage comes from the observation that the lateral dimension $W$ of a central structure is virtually independent of impact angle $\theta$ while crater diameter decreases (Schultz 1992c, 1994). Although both approaches to estimating impactor sizes agree well, the scatter among the data points is smaller when Eq. (15) is used (Fig. 10).

The next key variable is impact angle. The sector of missing uprange ejecta around an oblique crater is strongly influenced by impact angle (Gault and Wedekind 1978, Schultz 1992c). A lower angle impact results in a wider ejecta missing sector, and a higher angle impact results in a narrower ejecta missing sector. When the impact angle is less than about 40° (measured from the horizontal), the angle subtended by the sector of missing ejecta becomes a good index of impact angle. The estimation error of impact angle with this method is roughly 20% at lower angles and becomes less reliable at higher angles (Schultz 1992c).

Although impactor-size effects can be taken into account by normalization with impactor radius, the decreasing atmospheric density with height also affects the dynamics of an impact vapor cloud and atmospheric response for large impacts. To avoid this complication, we exclude craters larger than 70 km in diameter from the measurement. Small craters also can be problematic, principally due to atmospheric breakup of the impactor (Phillips et al. 1991). Thus we also exclude craters smaller than 20 km in diameter. Lastly, we have excluded craters on tesserae or other complex terrains that will affect the early-stage coupling process between projectiles and the target.

The result of the measurement clearly shows a correlation between the cotangent of impact angle ($\cot \theta$) and the downrange offset $L$ of a run-out-flow source region (Fig. 10). Although there is significant scatter in the plot for higher impact angles ($>30^\circ$), the normalized downrange offset $L/r_p$ can be fitted well with the relation

$$
\frac{L}{r_p} \propto \cot^\eta \theta, \quad \eta = 1 \pm 0.3.
$$

This relation should reflect three major processes: momentum transfer and energy transfer from an impactor to a downrange-moving vapor cloud and atmospheric interaction during downrange motion of the vapor cloud. The hydrocode calculations describe the atmospheric interaction and yield a relationship among $L/r_p$, $\phi_{energy}$, and $\phi_{momentum}$ as shown by Eq. (12). Thus if one of the partitioning mechanisms of momentum and energy is known, the other mechanism will be constrained by Eqs. (2) and (16). Laboratory experiments by Schultz and Gault (1990) indicate that the total momentum of the downrange-moving ricochet from an oblique impact is constrained by

$$
\phi_{momentum} \propto \cot^\xi \theta, \quad \xi \geq 0.7.
$$

The downrange-moving ricochet in laboratory experiments is not fully vaporized, but it reaches a vaporization condition easily in planetary impacts and becomes a vapor cloud owing to the much higher velocities. Equation (17) reflects momentum losses of the projectile during the penetration stage. In laboratory experiments, such losses largely reflect the deformation/fragmentation process of the projectile due to shock compression, spallation, and shear stresses developed during the penetration stage. Note that the relation (17) only holds over a certain range of impact angles (i.e., from about 10° to about 45° measured from the horizontal; Schultz and Gault 1990). Beyond this range, the momentum partitioning $\phi_{momentum}$ follows other trends, and hence it never exceeds unity at a very low angle $\theta$.

Because Eq. (17) is derived from laboratory-scale experiments, its direct application to planetary scales requires caution. Nevertheless, the geologic record on planetary surfaces strongly suggests that the momentum-coupling processes during a planetary-scale impact are similar to that observed in the laboratory. For example, many oblong craters on the Moon, Mars, and Venus have been found to be associated with downrange sibling craters (Schultz and Gault 1990, Schultz 1992c). The oblong shape of craters due to oblique impacts and associated sibling craters on the planetary surfaces closely resemble laboratory experiments where the initial collision results in the subsequent impacts of high-speed ricochets from the top of the impactor. Consequently, we use Eq. (17) to approximate momentum coupling during impact events on Venus.

We can now obtain a first-order approximation for the relation between energy partition and impact angle by combining the
results of hydrocode calculations (Eq. (12)), Venus observations (Eq. (16)), and laboratory experiments (Eq. (17)):

$$\phi_{\text{energy}} \propto \cot^{\chi} \theta, \quad \chi \geq 0. \quad (18)$$

Equation (18) indicates that more internal energy is partitioned into the downrange-moving vapor cloud for a shallow-angle than for a steep-angle impact. This result contrasts with a pure shock-heating mechanism, which is usually assumed in numerical calculations for impact-induced melting/vaporization.

Recent three-dimensional hydrocode calculation indicates that the internal energy partitioned to the projectile by pure shock heating during an oblique impact can be given by a simple power law (Pierazzo and Melosh 2000):

$$\phi_{\text{energy}} \propto \sin^{\chi} \theta, \quad \chi \sim 1.5. \quad (19)$$

This predicts that the ratio of the internal energy of the projectile component at $15^\circ$ to that at $40^\circ$ is only 0.26. Equation (18) based on Venus observation, however, indicates that the ratio is larger than unity, which is about 4 times the prediction from Eq. (19).

Even if the estimated error in the Venus observation is taken into account, the ratio does not go below 0.7. Consequently, the internal energy partitioned to the projectile component at low impact angles estimated in this study based on observations of run-out flows on Venus is significantly more than that estimated from pure shock heating. This comparison strongly suggests that there may be significant energy-partitioning processes other than the pure shock heating described by the Rankine–Hugoniot equations.

Recent laboratory experiments may provide clues for understanding the energy partitioning processes. Laboratory experiments by Schultz (1996) indicate that both mass and energy of impact-induced vapor clouds increase at lower impact angles (measured from the horizontal). This was explained by shear heating and melt/vapor generation by sibling impacts due to high-speed projectile fragments. Equation (18) obtained from Venus observations, therefore, may reflect similar heating processes during hypervelocity impacts at planetary scales. Spectroscopic observation of a jetting process by Sugita and Schultz (1999) also strongly suggests that shear heating may play an important role in impact vaporization processes.

Another possible process contributing to the enhanced energy partitioning to downrange-moving vapor clouds at lower impact angles, particularly at the extremely high atmospheric pressures on Venus, is the ablation of high-speed fragments and melt. Spectroscopic observations of atmospheric interactions with downrange impact vapor at high ambient pressures indicate that the high-temperature vapor emitting intense molecular radiation comes from vapor ablating from the surface of high-speed fragment/melt debris (Sugita and Schultz 1998). Ablation of high-speed fragments efficiently converts kinetic energy of ricocheted projectile fragments and downrange debris into internal energy (Schultz and Gault 1982). Because both the velocity of projectile ricochets and their kinetic energy are greater at lower impact angles, both the efficiency and the amount of converted kinetic energy of high-speed fragment/melt to internal energy of ablation vapor should also be greater at lower impact angles.

If the energy conversion takes place rapidly enough, the subsequent evolution of a vapor cloud augmented by the atmospheric ablation process is indistinguishable from that of a vapor cloud with initial internal energy created during the first compression stage.

### 4.2. Condensation Condition

Both laboratory experiments (Schultz 1996) and numerical calculations (O’Keefe and Ahrens 1986, Pierazzo and Melosh 2000) agree that the vapor cloud produced during a lower angle impact has lower temperature (specific internal energy) and higher translational velocity (kinetic specific energy). Although both energies contribute to the terminal average temperature of the vapor cloud in an atmosphere, the contribution of the initial internal energy is much more important (Eq. (13)). Consequently, vapor clouds produced by lower angle impacts are more likely to reach a condensation condition. This is consistent with the more frequent occurrence of run-out flows associated with craters due to lower angle impacts (Schultz 1992c, Chadwick and Schaber 1993).

It is noted that ablation of high-speed impact fragments discussed above may delay the condensation of impact vapor, particularly for lower angle impacts. Spectroscopic observations by Sugita and Schultz (1998) indicate that ablation vapor may have a temperature higher than that of impact vapor created during the initial compression stage. When such high-temperature ablation vapor is incorporated into a vapor cloud, the average specific internal energy increases, thereby delaying condensation. The effect of the ablation vapor is expected to increase for low-angle impacts. Laboratory experiments by Schultz and Gault (1990) indicate that the velocity of downrange-moving debris is higher at lower impact angles, thereby producing higher ablation temperatures (Sugita and Schultz 1998). Thus, ablation-induced vapor should delay the condensation at lower angles. The effect of a higher temperature ablation vapor at lower impact angles, however, may not be significant compared to that of the reduced temperature of vapor created during the initial compression stage. The mass of impact-induced vapor at planetary scales will be multiples of a projectile (e.g., O’Keefe and Ahrens 1977) even at low impact angles (Schultz 1996), but the mass of high-speed downrange fragments may only be a projectile mass at most (Schultz and Gault 1990). Since the mass of ablation vapor is less than or equal to that of high-speed fragments, changes in specific energy of ablation vapor as a function of impact angle may not affect drastically the average specific internal energy of the entire impact vapor cloud.

Because the hydrocode runs do not indicate large-scale mixing between impact vapor and surrounding air, the local total pressure in a vapor cloud is assumed to correspond to the partial pressure of impact vapor for estimating the condensation...
condition. The condensation process in the impact vapor, nevertheless, may be significantly complicated by mixing with ambient air. Such mixing most likely will prevent condensation or deposition of impact vapor and hence emplacement of run-out flows as discussed later in this section. Consequently, processes controlling mixing with air may provide further insights and constraints on the evolution of impact vapor in Venus's atmosphere. There are a couple of possible mixing processes.

First, if the Rayleigh–Taylor instability observed during the early stages continues to evolve, extensive mixing between impact vapor and ambient air may become significant. This process would delay condensation, since the temperature of the surrounding shock-heated air at this time is considerably higher than within the vapor (Figs. 3e and 3h). Atmospheric mixing also will decrease the partial pressure of the impact vapor component, which decreases its condensation temperature.

Second, mixing with the surrounding atmosphere also may occur during later stages, through buoyancy effects. As the vapor cloud decelerates, the heated cloud will rise due to buoyancy forces and entrain the surrounding air, thereby producing the classical vertical mushroom cloud. Note that this mushroom cloud formation is different from the mushroom-shaped dynamic disturbance observed in the hydrocode calculations (i.e., Figs. 3j, 3k, and 3m) but is a more classical one due to buoyancy forces (see, e.g., Jones and Kodis 1982). Because entrained air at this stage has not been heated by a strong shock, it has a much lower temperature than the rising impact vapor. This process should accelerate condensation of impact vapor. Condensation by this process, however, may not result in direct deposition. Because condensate particles are very small, settling rates for individual particles will be very slow. Significant deposition requires macroscopic gravitational collapse of a rising cloud similar to pyroclastic-flow formation. However, the extremely high temperature of the impact vapor is most likely capable of sustaining the buoyancy force to bring the cloud to high altitude in Venus's atmosphere just as a volcanic eruption cloud does (Sugita and Matsui 1993). Consequently, observed run-out flow emplacement indicates that impact vapor must have condensed before buoyancy forces lifted and mixed it with the cooler ambient air.

The requirement of rapid vapor condensation on Venus may provide further constraints on the nature of impactors. First, higher velocity impacts should generate higher temperature vapor at a given impact angle (e.g., Schultz 1996). Condensation and subsequent deposition of the resulting impact vapor is accordingly delayed and may be interrupted by the processes discussed above. Second, volatility of impactors affects the condensation condition of impact vapor. Because an increasing fraction of the downrange-moving vapor cloud is derived from the projectile for decreasing impact angles (Schultz and Gault 1990, Schultz 1996), downrange-moving vapor clouds created by oblique volatile-rich projectiles such as comets and carbonaceous asteroids are expected to contain a large amount of volatile material. Such volatile-rich vapor has lower condensation temperature, thereby reducing the chance of condensation. Consequently, both high impact velocities and volatile-rich impactors have less chance of forming run-out flows (Schultz 1992c).

Here it is noted that a general correlation is expected between volatility and impact velocity of projectiles. Volatile-rich asteroids (i.e., C and D types) and comets have larger heliocentric distances than asteroids with more silicate and metal (i.e., S and M types) (Gradie and Tedesco 1982). Those objects with larger heliocentric distance will have higher orbital velocities relative to Venus. The combined effect of the velocity and the volatility of impactors may account for exceptions to the more frequent occurrence of run-out flows around asymmetric craters that are inferred to be created by oblique impacts (<45°) on Venus. If metal/silicate-rich asteroids collide with Venus at low velocities, condensation of resulting impact vapor clouds is very likely. Vapor clouds due to such impacts will condense even when the impact angle is relatively high. This may account for exceptional run-out flows around craters with high axisymmetry. Conversely, if volatile-rich asteroids or comets collide with Venus at high velocities and low angles, they may result in asymmetric craters without run-out flows or with diffuse flows, as observed on Venus (Schultz 1992c).

4.3. Wind Generation

Another important observable consequence predicted from the hydrocode calculation is strong wind generation by both the early-stage shock wave and the late-stage vortices around an impact vapor cloud (Figs. 3l and 3o). If the radius of the projectile that induced the impact vapor cloud were the same as that of the cloud (i.e., 1 km), the final crater size predicted from a crater scaling law by Schultz (1992c) based on results by Schmidt and Housen (1987) and Gault and Heitowit (1963) would be about 8 km for a 30° impact angle measured from the horizontal. If a downrange-moving vapor cloud contains a significant amount of target material (i.e., \( M_{\text{vap}} / M_{\text{proj}} > 1 \)), the final crater radius will be smaller than this value. Thus, a large fraction of some of the vortices observed in the later stage of the calculation (Figs. 3l and 3o) are more than a crater radius away from the impact center, and some are more than two crater radii away. Also, the initial outward blast wave has traveled several crater radii by this point. Consequently, winds generated by vapor cloud motion may interact with the ground surface outside both the crater cavity and the ejecta deposit.

In fact, fresh craters on Venus exhibit complex wind-streak patterns that appear to be controlled by impact-generated processes (Schultz 1992c). The wind patterns documented in the hydrocode calculations provide important clues for understanding these enigmatic wind patterns around craters on Venus that exhibit outward, inward, and circumferential patterns. First, the ground surface is scoured by the intense primary shock-induced wind (Figs. 3c, 3f, and 3i). Second, this initial outward/downrange wind is followed by an intense inward/uprange wind (Fig. 3i). The sharp contrast in temperature (Fig. 3h) and a large velocity jump (Fig. 3i) indicate that this wind is also a shock wave. Consequently, dislodged and mobilized material...
can be quite coarse (i.e., radar bright), as is observed. Third, the primary shock wave detaches itself from the vapor cloud and dissipates quickly. The secondary shock wave, however, results in several vortices that continue to develop as the primary shock waves dissipate (Figs. 3l and 3o). These vortices will leave the last surface features on Venus among the vapor-induced phenomena. Because these vortices still generate winds (~500 m/s) sufficient to mobilize large particles (~1 m; Greeley and Iversen 1984), their surface expression also should be radar bright. Nevertheless, their lower intensity is insufficient to erase shock-driven outward/inward wind streaks created earlier. Air motion due to such late-stage vortices may also interact with regional/global winds in Venus's atmosphere and induce a circumferential wind pattern around craters (Schultz 1992c). Consequently, complex cross-pattern wind streaks associated with a fresh crater on Venus can be created by the atmospheric response to a time-evolving impact vapor cloud.

5. CONCLUSIONS

Two-dimensional hydrocode calculations were performed to understand the dynamics of expanding vapor clouds with translational motion within an ambient atmosphere. The initial conditions such as temperature, density, and velocity expected for downrange-moving vapor clouds due to planetary-scale oblique impacts were used. Calculations revealed several intriguing hydrodynamic phenomena:

1. An intense shock front forms in front of a downrange-moving impact vapor cloud, behind which most of the cloud’s vapor mass is gathered and forms a hemispherical shell.

2. The intense frontal shock controls the shape of an impact vapor cloud and results in dynamic reshaping. Such reshaping decreases the deceleration rate of impact vapor clouds because it reduces the air drag force. This process is particularly important during early stages when shock intensity is high.

3. Rayleigh–Taylor instabilities develop on the boundary between the shell of the impact vapor and the shocked air into which it travels.

4. After the initial frontal shock is detached from the impact vapor cloud, a strong secondary shock wave propagates backward through the cloud and reheats the impact vapor.

5. The backward shock wave results in a mushroom-shaped disturbance behind the impact vapor cloud.

6. As a result of shock wave motion induced by an impact vapor cloud, strong and large vortices form and last long after shock waves dissipate.

The center of mass and the average temperature of the downrange-moving impact vapor cloud were monitored throughout each computational run. The time history of these quantities revealed a number of important aspects of dynamic evolution of an impact vapor cloud:

1. Expansion of an impact vapor cloud increases its projected surface area, thereby increasing air drag forces very rapidly. This results in very efficient deceleration of a vapor cloud at late stages.

2. After rapid deceleration, an impact vapor cloud does not stop but oscillates back and forth. This oscillation lasts long after the initial shock waves have dissipated.

The time history also indicates that the vapor-cloud-origin hypothesis proposed for run-out flows around Venus craters satisfies several key requirements quantitatively:

1. The downrange-moving impact vapor cloud travels beyond the final crater rim well before crater growth and ejecta emplacement.

2. The terminal average temperature of an impact vapor cloud is close to the condensation temperature of a typical geologic material.

3. The cooling time scale for impact vapor is shorter than that for ejecta emplacement.

In order to understand what controls the total travel distance of an impact vapor cloud and the condensation of impact vapor, we performed a series of parametric studies of the initial conditions. The results of the parametric study indicate the following:

1. As a result of reduced air drag force due to dynamic reshaping, the rate of increase in the total travel distance of downrange-moving impact vapor cloud for increased initial translational velocity is higher than a linear rate.

2. The travel distance normalized by the impactor radius depends little on the mass of an impact vapor cloud but significantly on both energy and momentum partitioned to the vapor cloud.

3. The terminal temperature of an impact vapor cloud increases linearly with the effective initial specific energy of the cloud, which is the sum of both the specific internal energy and 20% of the specific kinetic energy.

4. Calculated terminal temperatures for vapor clouds with initial conditions expected for Venus impacts range from below to above the condensation temperature depending on impact velocity and vapor composition.

The above calculation results have significant implications for both geology on Venus and impact mechanics at planetary scales:

1. The downrange offset of run-out-flow source regions around asymmetric craters (i.e., oblique impacts) on Venus suggests that energy partitioned into a downrange-moving vapor cloud is higher at shallower impact angles.

2. The inferred enhancement in internal energy partitioned to a downrange-moving vapor cloud at shallow impact angles is not readily explained by a pure shock-heating mechanism but is consistent with both a shear heating mechanism during an impact and secondary vaporization due to ablation of high-speed fragment/melt ricocheted from an impactor.

3. An impact vapor cloud with a lower initial temperature, which is expected for a shallower angle impact, is more likely to
reach a condensation condition. This may account for the more frequent occurrence of run-out flows associated with craters due to oblique impacts on Venus.

(4) Both impact velocity and projectile composition control the condensation condition of the downdrange component of an impact vapor cloud. This may account for exceptions to the relation observed on Venus between occurrence of run-out flows and inferred impact angles. Volatile-rich projectiles at high velocities may result in craters without run-out flows even at low impact angles, whereas silicate/metal-rich projectiles at low velocities may yield run-out flows around symmetric craters, which are inferred to be produced at high impact angles.

(5) Superimposed wind patterns created during different evolutionary stages of a downrange-moving impact vapor cloud may account for complex patterns of wind streaks around fresh impact craters on Venus.

**APPENDIX**

**Time Scale of Radiative Cooling**

When a high-temperature gas body is put into a vacuum or an optically transparent medium with a low temperature, it radiates and cools. To obtain an exact solution of this problem, one must solve the radiative transfer equation. This equation is not only mathematically complex but also requires knowledge of the opacity of the gas as a function of both temperature and pressure. Fortunately, however, there are simple approximate solutions for special conditions that allow constraining the consequences. Using such solutions, we can assess the radiative cooling rate of impact vapor clouds within Venus’s atmosphere.

If the optical thickness of a gas body is sufficiently larger than unity over the spectral range of its thermal radiation, the radiant flux density of the gas body is given by Stefan–Boltzmann’s law. If the shape of the gas body is spherical, its temperature change is described by

\[ \frac{4\pi}{3} R^3 \rho C_p \frac{dT}{dt} = -4\pi R^2 \sigma T^4, \]  

(A1)

where \( T, T_i, R, \rho, C_p, \) and \( \sigma \) are mean temperature, surface temperature, radius, mean density, heat capacity, and the Stefan–Boltzmann constant, respectively. Because the radiative energy flux from the body is proportional to its surface area, this process is called as “surface cooling” (e.g., Zel’dovich and Raizer 1967). Then the time scale for this cooling process is

\[ \Delta t_{\text{ref}} \geq \frac{\beta C_p R}{3\sigma T_i^3}, \]  

(A2)

where \( T_f \) is the final temperature, and the initial temperature is assumed to be much higher than \( T_f \). The equality holds when the surface temperature is always equal to the mean temperature. Converting this time scale (A2) into a dimension of velocity, we obtain

\[ v_{\text{ref}} \equiv \frac{R}{\Delta t_{\text{ref}}} \leq \frac{3\sigma T_f^3}{\beta C_p} = 10 \, \text{m/s} \times \left( \frac{T_f}{10,000 \, \text{K}} \right)^3 \left( \frac{\beta}{50 \, \text{kg/m}^3} \right)^{-1}. \]  

(A3)

This quantity can be compared with the sound velocity of the gas if adiabatic expansion of a gas is a competing cooling mechanism, because the rarefaction wave, which controls adiabatic expansion, travels at the local sound speed. The sound velocity of an impact vapor cloud due to a hypervelocity (>15 km/s) collision is initially several kilometers per second and eventually decreases to several hundred meters per second. Thus the rate of radiative cooling is an order of magnitude (or more) lower than that of adiabatic cooling in an expanding impact vapor cloud in Venus’s atmosphere. Note that the cooling rate (i.e., Eq. (A3)) is independent of opacity; it depends only on the gas temperature and density.

When the optical thickness of a gas sphere is not large enough, the total radiative output from the body cannot be expressed in a form as simple as Eq. (A1). The total radiative output from the body, nevertheless, cannot exceed the quantity predicted by Eq. (A1) if the temperature in the body is uniform (see, e.g., Cannon 1985). Thus when a vapor cloud has a homogeneous thermal structure, the rate of radiative cooling is much lower than the adiabatic expansion indicated by Eq. (A3). If a gas body has a strong thermal anomaly inside, a similar approach can be applied by neglecting incoming radiation from the surrounding lower temperature part of the body and considering the high-temperature portion as an effectively isolated gas body. Then the radiative cooling is again constrained by Eq. (A3).

It should be noted, however, that an extremely high temperature such as 10^8 K actually brings the radiative cooling speed (Eq. (A3)) to 1 km/s. Such temperatures may occur only at the very early stage in the evolution of the impact vapor cloud. In this stage, Eq. (A3) no longer holds because it neglects the effects of radiative cooling. However, because the opacity of a gas generally becomes very large at such high temperatures due to dissociation and ionization, it becomes effective to follow another approach using a diffusion approximation (e.g., Zel’dovich and Raizer 1967),

\[ \rho C_p \frac{dT}{dt} = -\nabla \left( \frac{16}{3} \sigma T^4 \nabla T \right), \]  

(A4)

where \( l \) is the mean free path of photons. Then the time scale \( \Delta t_{\text{rad}} \) of the radiative diffusion over the length of \( L \) is given by

\[ \Delta t_{\text{rad}} = \frac{3 \rho C_p \rho^2 L^2}{16 \left( \frac{\beta}{\pi} \right)^2}, \]  

(A5)

where \( k \) is the extinction coefficient of photons and is related to the mean free path of photons through

\[ l = \frac{1}{k \rho}. \]  

(A6)

Experiments of heated atmospheric air indicate that \( k \) is about 0.1 m^2/kg at a temperature of 10,000 K and a density of 1 kg/m^3 (Churchill et al. 1966, Zel’dovich and Raizer 1967). This value is clearly not for the very early stage of impact vapor clouds. The chemical composition is very different, and both temperature and density are lower. However, the gases of interest, such as silicate vapor and carbon dioxide, are expected to have a larger extinction coefficient due to their more complex molecular structure than diatomic terrestrial air. Moreover, both temperature and density increase the extinction coefficient. Consequently, the use of the above value of \( k \) gives a lower limit for the radiative time scale. The time scale for radiative cooling \( \Delta t_{\text{rad}} \) for a thermal anomaly with wavelength \( L \) of 100 m, temperature \( T \) of 10^6 K and density \( \rho \) of 100 kg/m^3 is about 10 hours. Similarly, \( \Delta t_{\text{rad}} \) is about 1.7 min for \( L = 100 \) m, \( T = 10^5 \) K, and \( \rho = 100 \) kg/m^3, and it is 1.4 min for \( L = 100 \) m, \( T = 10^4 \) K, and \( \rho = 5 \) kg/m^3. These time scales are much longer than the duration of such high temperature (Fig. 5). Consequently, radiative cooling can be neglected even in the very early stage of the vapor cloud expansion.

**ACKNOWLEDGMENTS**

The authors thank T. Yabe and T. Aoki for helpful comments on numerical methods. This study also benefited from constructive discussions with D. A. Crawford and O. S. Barnouin-Jha. This research was supported by NASA Grant NAGW-705.
REFERENCES


