Notes

Observations on the Effects of Materials on Card Gap Test Results

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THE dependence of the card gap value for a propellant on the container material has been previously reported.[1] The data reported here on nitromethane (NM) and a mixture of tetranitromethane (TNM) and acetonitrile (AN) confirm this dependence and, moreover, uncover certain effects that could be interpreted in terms of specific chemical or catalytic activity.

This card gap test is a standard test in the rocket industry used to determine the relative sensitivities of propellants to hydrodynamic shock.[2] It is capable of good reproducibility if one conforms to the conditions specified for the test. However, if the conditions are varied, then marked changes in the card gap value for a given propellant are possible.[1] The simplicity and reproducibility of the test and its sensitivity to experimental conditions suggested that it could be useful in obtaining qualitative or semiquantitative information on the effects of certain variables on the initiation and detonation processes of monopropellants. It was therefore decided to use this method to study two monopropellants. The variable of greatest interest was the propellant container material. In addition to varying container materials, small quantities of several solids were added to the monopropellants to assess the presence or absence of particular chemical or catalytic effects.

The basic experimental procedure was that of the JANAF Panel on Liquid Propellant Test Methods.[2] The propellants were reagent grade NM and a mixture of TNM (99 + % by weight) and reagent grade AN in the ratio of 2 moles of AN to 1 mole of TNM.

The container materials investigated were Teflon-coated steel (the standard container material), stainless steel 304, aluminum (T-3003), Pyrex glass, and Lucite. These materials were selected because of large differences in their chemical and mechanical properties. The dimensions of the cups prepared from these materials were very close to the dimensions of the standard cup. In some tests, small quantities of a variety of different solids, usually in the form of spheres (less than 1 cm^3 volume, equivalent to 7 to 10 in number), were added to the test containers. The card gap value was located to ±5 cards, since only gross changes as a function of the variables were sought.

The card gap values of NM in the standard Teflon-coated steel, stainless steel, aluminum, Pyrex glass, and Lucite cups were 15–20, 20–30, 20–30, < 5, and < 5, respectively. The addition of solids (carbon steel, aluminum, glass, sand, and stones) to NM in the stainless steel and aluminum cups raised the card gap value by at least 10 cards to between 40 and 50.

The card gap values for TNM-AN in the standard Teflon-coated steel cup, stainless steel, and aluminum were 75, 95–100, and 180–200, respectively. Carbon steel, glass, and aluminum were added to TNM-AN contained in the stainless steel cup and resulted in an increase of card gap value from 95–100 to 145–165, 110–120, and to > 230, respectively.

In considering the effects of variations in container material and added solids on the card gap value of NM, no effect that can be uniquely attributed to the chemical nature of the material in contact with the NM can be discerned. The added solids tend to raise the card gap value, probably as a result of the enhancement of "hot-spot" formation in the detonation initiation process,[3] but all of the added solids are about equally effective. The card gap values of NM in the nonmetallic containers were considerably less than those in the metallic containers. The reasons for this are not precisely known at this time, nor can the fairly small differences between the card gap value of NM contained in Teflon-coated steel, and stainless steel or aluminum be explained. The tests are not sufficiently quantitative to justify a correlation with variations in chemical, physical, or mechanical properties of the materials in question.

The studies with TNM-AN were more limited than those with NM in that no nonmetallic containers were studied. Also, because TNM-AN is considerably more sensitive than NM, caution must be exercised in the interpretation of the results, since it has been shown that the initiating pressure pulse passing through a stack of cards falls off exponentially with increasing number of cards.[4] This notwithstanding, it appears that the marked effect of aluminum, whether it is in contact with TNM-AN as the container material or as an additive, in increasing the card gap value is evidence of some specific chemical or catalytic effect on the initiation process.

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Influence of Conduction on Spacecraft Skin Temperatures

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Nomenclature

A = area of shell
k = thermal conductivity of shell material
r = radius of shell
t = shell thickness
T = absolute temperature

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$\bar{T}$ = reference temperature, Eq. (2)
α = spacecraft skin conductance parameter, $kt/(4r^2\sigma\bar{\epsilon}\bar{T}^3)$, dimensionless
$\bar{\alpha}$ = over-all absorptivity for solar radiation
β = relative internal radiation parameter, $\epsilon_i/\bar{\epsilon}$
$\delta(\theta)$ = operator: $\delta(\theta) = \begin{cases} 1, & |\theta| \leq \pi/2 \text{ (sun)} \\ 0, & |\theta| > \pi/2 \text{ (shadow)} \end{cases}$
$\bar{\epsilon}$ = over-all infrared emissivity of external surface
ϵ_i = infrared emissivity of the interior
θ = polar coordinate, $0 \leq \theta \leq \pi$; also the angle between the normal and the direction of sun rays
σ = Stefan-Boltzmann constant
τ = dimensionless temperature
τ_0 = dimensionless temperature, no conduction

Introduction

THE question of temperature distribution on the surface of space vehicles has already received some attention in the literature[1–4]; in the present note, expressions are derived for local temperature distribution over the surface of a thin-walled spherical spacecraft for the situations where the fourth-power law of radiation must be preserved. The local temperature distribution over the spacecraft surface and, in particular, the maximum skin temperature are of importance primarily with respect to the performance of solar cells that ordinarily cover a considerable portion of its external area.

Assuming that a good portion of the spacecraft interior is empty, it is advisable to include in the present analysis also the effects of internal radiation.

Statement of the Problem

The energy balance for an element of a thin spherical shell, permanently exposed to solar radiation and radiating to space at absolute zero, yields the relation

$$\frac{kt}{r^2}\left(\frac{d^2T}{d\theta^2} + \cot\theta\frac{dT}{d\theta}\right) + \delta(\theta)S\bar{\alpha}\cos\theta + \frac{1}{4}\frac{\epsilon_i}{\bar{\epsilon}}S\bar{\alpha} - \sigma(\epsilon_i + \bar{\epsilon})T^4 = 0 \quad (1)$$

Since the temperature distribution must be symmetric, $\partial T(0)/\partial\theta = \partial T(\pi)/\partial\theta = 0$.

Equation (1) is the familiar Laplace equation in spherical coordinates with θ as the only variable, with three additional terms that are due to the effects of radiation. It is assumed here that $t \ll r$, and there is no temperature drop through the thickness of the shell. The radiation terms are, in their respective order, 1) direct solar radiation, 2) radiation absorbed internally, and 3) the energy lost internally and externally according to the Stefan-Boltzmann law. The internal radiation term is obtained from integrating the contribution of all area elements of the shell.

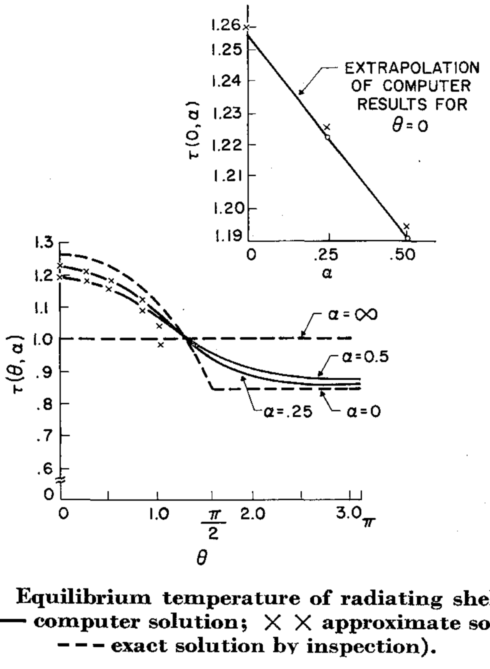

Fig. 1 Equilibrium temperature of radiating shell (β = 1.0; —— computer solution; × × approximate solution; - - - exact solution by inspection).

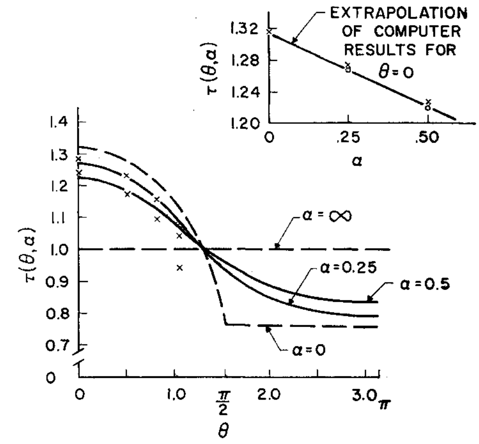

Fig. 2 Equilibrium temperature of radiating shell (β = 0.5; —— computer solution; × × approximate solution; - - - exact solution by inspection).

It is convenient to reduce Eq. (1) to a dimensionless form, using as the reference temperature $\bar{T}$, defined as

$$\bar{T} = [S\bar{\alpha}/4\sigma\bar{\epsilon}]^{1/4} \quad (2)$$

This is the equilibrium radiation temperature of a sphere exposed to solar flux for the limiting condition of either $k \to \infty$ or $\epsilon_i/\bar{\epsilon} \to \infty$. For a thin shell with a finite k and β, it may be written

$$\bar{T} = \left[\frac{1}{A}\int_A T^4 dA\right]^{1/4} \quad (3)$$

Equation (1) may be rewritten in dimensionless form as

$$(\tau'' + \cot\theta\tau')\alpha + \delta(\theta)\cos\theta + \frac{\beta}{4} - \frac{1+\beta}{4}\tau^4 = 0 \quad (4)$$

with $\tau'(0) = \tau'(\pi) = 0$.

Equation (4) represents a nonlinear boundary value problem, where the difficulty lies in getting the value of the maximum temperature $\tau(0)$. For an α large enough so that $|\tau - 1|/\tau \ll 1$, Eq. (4) may be linearized; then, solution in terms of Legendre polynomials is possible.[4]

For very thin shells, of particular importance is the situation where α remains so small that the effects of shell conductance introduce only a very small deviation from the temperatures based on radiative equilibrium alone. In this case α can be set equal to zero, and the remaining terms of Eq. (4) may be solved for τ to give τ_0:

$$\tau_0 = 2^{1/2}\left[\frac{[\delta(\theta)\cos\theta + (\beta/4)]}{(1+\beta)}\right]^{1/4} \quad (5)$$

which represents here the "zero-order" approximation to the solution. For $\alpha \neq 0$, approximate analytical methods must be used. Let us assume, therefore, that, for a sufficiently small value of α and for a suitable region $0 \leq \theta < \theta_0$, one can write a two-term approximation to the exact solution of Eq. (4)† as

$$\tau = \tau_0 + \tau_1\alpha + \ldots \quad (6)$$

By a direct substitution of this expression in Eq. (4) and

† It is readily possible to introduce terms of the order of α^2 and higher by this method, but for $\beta \neq 0$ a good degree of accuracy is available already with the two-term approximation.

Table 1 Comparison of results of numerical integration of Eq. (6) with two-term approximation with internal radiation absent

$\cos\theta$	τ, Ref. 4	τ, Eq. (6)
1.000	1.335	1.352
0.875	1.281	1.293
0.750	1.226	1.223
0.500	1.108	1.002

by collecting the terms according to powers of α, one obtains for τ_1

$$\tau_1 = -(1 + \beta)^{-2}\tau_0^{-6}[2\cos\theta + 3\sin^2\theta(1 + \beta)^{-1}\tau_0^{-4}] \quad (7)$$

Solutions using Eq. (7) are plotted in Figs. 1 and 2. As may be noted from the figures, temperature curves with α as a parameter cross in the region $\pi/3 < \theta < \pi/2$. The temperature, therefore, ceases to be an analytic function of α in that region, which in effect restricts the range of validity of Eq. (6) to about $0 \leq \theta < \pi/3$.

It is interesting to note that all temperature lines seem to cross in the point where $\tau \approx 1$, in the neighborhood of an inflection point. This fact can be used to get information about the temperatures where Eq. (6) breaks down. From the comparison with the data in the Figs. 1 and 2 and those by Nichols,[4] it is seen that the value of the abscissa of the point $\tau = 1$ is close to $\theta = \cos^{-1}\frac{1}{4}$ for the range of α under consideration. Then, using Eq. (6) for $\theta < \pi/3$, together with the information on the approximate location of the point $\tau = 1$, the temperatures over the rest of the shell can be calculated, using as a check the relation implied in Eq. (3).

It is also interesting to discuss the way in which Eq. (6) works in the case of negligible internal radiation ($\beta = 0$), where the deviation of the shell temperature from the no-conduction condition becomes relatively more significant. In Table 1, the data from Ref. 4, p. 28, are shown for comparison with the figures based on $\alpha = 0.25$ and $\beta = 0$.

From Table 1 the conclusion can be drawn that, for $\beta = 0$, Eq. (6) still gives a reasonable approximation, showing to what extent conduction effects alone will reduce the adiabatic wall temperature on a thin spherical shell in space (typical for $0 < \alpha < 0.5$).

Concluding Remarks

The temperatures obtained by the methods discussed in the foregoing are useful in the calculations connected with the output of the spacecraft solar cell powerplants.

However, one more interesting application seems to be possible. As suggested by Reismann and Jurney,[5] the energy flux F for a sphere near the stagnation point may be approximated for hypersonic speed by the equation

$$F = a + b\cos\theta$$

a and b being functions of the freestream Mach number. They may be considered to be constants for conditions where the effects of radiative cooling are of importance. From an inspection of Eq. (1), it is obvious that it is also applicable for this situation, with only minor modifications, if the freestream velocity remains parallel to the ray $\theta = 0$.

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Satellite Dynamics for Small Eccentricity Including Drag and Thrust

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Nomenclature

A_0	=	satellite's cross-sectional area
C_D	=	drag coefficient
F_r, F_v	=	thrust component, radial and normal, respectively
m	=	mass of satellite
μ	=	gravitational force on unit mass at unit distance
V	=	velocity of satellite
ρ	=	density of atmosphere
δ	=	$\frac{9}{2}(K + \frac{3}{2})^2 - \frac{1}{2}$

THE motion of a satellite in an orbit of small eccentricity was variously considered in the literature, notably by Perkins,[1] Lawden,[2] Newton,[3] Karrenberg et al.,[4] and Parsons.[5] The nonlinear problem with exponential atmosphere has been the subject of outstanding papers by King-Hele,[6] Sterne,[7] and others. In Ref. 1 the linearized problem is treated for constant thrust without consideration of drag; the results are used in Ref. 4 where the drag is treated as a negative (constant) thrust assuming an atmosphere of constant density and assuming that the initial orbit is *exactly* circular. A complete linearized solution, that is, one valid to the first order in $(\Delta r/r)$, including *arbitrary* initial conditions, *variable* density and drag, and an *arbitrary* thrust program is not known to this author. The availability of explicit expressions in closed form is quite useful in preliminary orbit determination (for instance, for a two-point boundary value problem), in rendezvous problems, and in low thrust problems; hence, the necessary calculations were carried out,[8] and the conclusions are presented in the following note.

For a spherical planet with stationary atmosphere, the equations of motion in polar coordinates (r and v) are

$$m(\ddot{r} - r\dot{v}^2 + \mu/r^2) = F_r(t) - \tfrac{1}{2}C_D A_0 \dot{r} V \rho(r) \quad (1)$$

$$m[(r\ddot{v}) + 2\dot{r}\dot{v}] = F_v(t) - \tfrac{1}{2}C_D A_0 r\dot{v} V \rho(r) \quad (2)$$

Let the satellite be observed, at time $t = 0$, to have angular velocity $\dot{v}(0)$ at distance r_0; an angular velocity n_0 $[\neq \dot{v}(0)]$ is defined by the well-known relation

$$r_0 n_0^2 = \mu/r_0^2 \quad (3)$$

and the unknown functions $\epsilon(t)$ and $\varphi(t)$ are introduced

$$r(t) = r_0[1 + \epsilon(t)] \quad (4)$$

$$\dot{v}(t) = n_0[1 + \varphi(t)] \quad (5)$$

into the basic equations, which are then linearized by neglecting higher powers of ϵ and φ. For instance, the velocity $V = (\dot{r}^2 + r^2\dot{v}^2)^{1/2}$ becomes $V = r_0 n_0(1 + \epsilon + \varphi)$ to this approximation. Similarly, the linear approximation to the exponential atmosphere $\rho = \rho_0 \exp[(r_0 - r)/H]$ for scale height H becomes

$$\rho(r) = \rho_0(1 - K\epsilon) \quad (6)$$

where $K = r_0/H$ is introduced as a dimensionless quantity. The initial conditions on ϵ and φ at $t = 0$ are

$$\epsilon(0) = 0 \qquad \dot{\epsilon}(0) = An_0 \equiv \dot{r}(0)/r_0 \quad A \ll 1 \quad (7)$$

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