Threshold-Determining Mechanisms for Discharges in High-Voltage Solar Arrays

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A theory is developed to account for the observed properties of discharges of solar arrays immersed in plasma. The theory is based on the assumption that a thin layer of insulating contaminant covers any metallic surface exposed to the plasma. Ions of the space plasma, attracted by the negative potential on the array, neutralize on the layer's surface, resulting in a buildup of electric field in the layer. With continued charging, the internal field becomes sufficiently large to cause emission across the metal/layer interface and subsequent ionization and electron heating within the layer. The hot electrons can be emitted into vacuum; this emission constitutes the positive feedback mechanism leading to discharge. A mathematical description of the processes based on the foregoing assumptions is developed. Quantities derivable from the theory include a voltage threshold for arcing. The existence of the threshold and its predicted weak dependence on plasma density appear consistent with experimental results.

Nomenclature

A $= 1.5 \times 10^6 \times 10^{4.52\phi^{-1/2}}$ amp/cm^2
B $= 30\phi^{3/2}$ (eV)MV/cm
d = effective ionization thickness of insulator
D_+ = hole diffusion coefficient
E = electric field
E_c = a property of insulator, having dimensions of electric field, characterizing electron emission from insulator
E_e = insulator band gap
E_o = vacuum electric field at insulator surface
E_i = electric field at metal/insulator interface
E_{ot} = threshold value of E_o
E_{it} = threshold value of E_i
H = property of insulator, with dimension of electric field, characterizing dependence of collisional ionization on electric field
I = total current
j_e = electron emission current
j_{FN} = Fowler-Nordheim emission from metal into conduction band of insulator
j_i = ion current density impinging on insulator surface
j_1 $= j_i + j_e$
j_2 $= j_{FN}$
m $= p(E_o)/p(E_{ot})$
m_e = electron mass
n_e = density of conduction electrons in insulator
n_+ = hole density
p = probability that a conduction electron at insulator/vacuum interface is emitted into vacuum
R = arc rate
t = time
T $= R^{-1}$
v = hole drift velocity
V = voltage
W = effective work function for electron emission from insulator's conduction band into vacuum
x $= E_i/B$
α = ion pairs produced by an electron per cm of path in insulator
α_o = high field limit for ion production by collisions in a unit length of path
ϵ = mean fractional energy loss per collision
ϵ_o = permittivity of vacuum
ϕ = metal/insulator work function
θ = electron temperature (in energy units) in insulator
θ_o = lattice temperature
K = dielectric constant of insulator
μ = electron mobility
μ_+ = hole mobility
τ = mean effective collision time for electrons in insulator

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I. Introduction

SNYDER[1,2] and Miller[3] have conducted experiments on the arcing of negatively biased solar arrays with interconnects exposed to plasma. Among the principal observations were

1) The rate of arcing increases, apparently in direct proportion, with the ion density.

2) The arcing rate is a sensitive function of voltage, exhibiting something resembling threshold behavior in the neighborhood of −200 V.

3) Before discharge, coverslip potentials swing toward negative values, indicating electron buildup on the coverslip as a precursor to discharge.

Jongeward et al.[4] have suggested that arcing is initiated as a result of ion neutralization and associated charge buildup on a thin insulating layer over the metallic interconnect. Indeed, the order of magnitude of arc rate in some experiments is consistent with the time required for ions attracted from the plasma to build up electric fields to the 10^7 volt/cm level on such an insulating layer. Moreover, the calculations of Jongeward et al. showed that the negative charge reaching the coverslip prior to arcing could not have come from the plasma but was instead the result of electron emission from the interconnect. There is, in fact, a substantial body of data in support of the general idea that nonmetallic emission processes are largely responsible for the ease, relative to Fowler-Nordheim tunneling, with which nominally metallic surfaces emit electrons.[5]

In general, the arc rate

$$R = f\ [j, V, (\alpha)] \tag{1}$$

may depend not only on current density j but on bias voltage V and other parameters as well; α represents the set of all other parameters, including the physical and geometric properties of the insulator.

Consider a simple insulator, one having a filled valence band and an empty conduction band, and ignore for the moment any traps that may exist within the band gap. Suppose that electrons are injected into the conduction band of the insulator, for example by field emission from an adjacent metal. Upon gaining energy from the electric field and drifting toward the insulator's surface, these conduction electrons and any others that approach the vacuum surface from within the insulator will be emitted into vacuum at the rate j_e (cm^{-2}) over a barrier with an effective work function W substantially smaller than the barrier $E_e + W$ required to liberate electrons from the valence band; if the conduction electrons are thermalized by some process, the emission will be thermonic in nature:

$$j_e = n_e (\theta / 2\pi m_e)^{1/2}\ e^{-W/\theta} \tag{2}$$

II. Electron Temperature

A complete mathematical description of the reaction kinetics requires a rate equation for n_e and an energy balance equation for for θ, assuming that the population n_e can be characterized by a temperature. Such a system of equations could easily be developed following lines pioneered by Rose[6] and Lampert.[7] The results would, in any event, be inconclusive, considering our lack of knowledge of the specific insulator we are dealing with and the wide range of values of the insulator's physically relevant parameters. Therefore, we shall keep our mathematical description as simple as possible, our main objective being to exhibit what may be the phenomenology of prebreakdown emission currents rather than to give an elaborate description of the underlying physics.

To begin, fix n_e and consider an equation for electron temperature θ. We suppose that θ is determined by a balance between energy gain from the electric field and loss to the lattice or to electrons in trapped states between valence and conduction bands. The processes by which a conduction electron loses energy to the lattice undoubtedly involve phonon emission, a very inefficient process; efficient thermalization of conduction band electrons probably occurs through their interaction with electrons in shallow traps located below the bottom of the conduction band. The energy balance between conduction electrons and the insulator is described by

$$\frac{d\theta}{dt} \approx \mu E^2 - \frac{\epsilon}{\tau}(\theta - \theta_o) \tag{3}$$

or, assuming equilibrium between heating and cooling rates,

$$\theta = \theta_o + \frac{\tau\mu}{\epsilon} E^2 \tag{4}$$

For $\tau\mu/\epsilon = 10^{-11}$ ($\tau \sim 10^{-14}$ s, $\epsilon = 0.1$, $\mu = 100$ cm^2/V-s), a reasonable value for many insulators, $\theta = 0.1$ eV at $E = 10^5$ V/cm. Equation (4) exhibits the sensitivity of the temperature to the properties of the insulator. Moreover, $\tau\mu/\epsilon$ may depend on E. In any case, it is noteworthy that the straight-line emission characteristic given in Latham[6] (Fig. 4, log I/V^2 vs $1/V$) can also be fitted rather well as a straight line in many other ways, including on a log I/V^2 vs $1/V^2$ diagram. It would appear then that the hot electron model of the type considered here can lead to the voltage dependence of arc rate observed in the NASA/LeRC experiments. To be in reasonable accord with these observations, however, the field E that determines the onset of significant emission must be the vacuum field at the surface of the dielectric as opposed to the possibly much larger field in the interior of the dielectric beyond the charge deposition layer near the surface.

The charge deposition layer consists of hot electrons and an excess of positive charge represented by holes in the valence band of the insulator. Both components of the charge migrate under the influence of the electric field and electron and hole density gradients. In the absence of collisions, hot electrons (0.1 eV) may migrate up to 10^{-5} cm away from the surface into the interior of the insulator before having their motion arrested by the electric field. The positive holes acquire drift velocities

$$v = \mu_+ E + D_+ \frac{\partial n_+}{\partial x} \tag{5}$$

of unknown magnitude. However, for $E \sim 10^5$ V/cm, μ_+ as small as 10^{-11} cm^2/V-s leads to hole migration distances of $\sim 10^{-6}$ cm in one second. The principal point is that the positive charge is unlikely to remain frozen at the surface, and for hole migration over distances $\geq 10^{-7} - 10^{-6}$ cm between discharges it is reasonable to suppose that the emission of any electrons that approach the insulator-vacuum interface is controlled by the vacuum electric field at the insulator surface.

III. Enhanced Electron Emission

In order for the enhanced electron emission to significantly increase the arc rate over that which would occur in the absence of emission, that is, with surface charging only by incident ions, it is necessary that the emission rate exceed the ion arrival rate.

To see how this may occur, consider a simple model in which the primary mechanism for enhanced electron emission is taken to be collisional ionization produced by field acceleration of electrons emitted from the metal into the conduction band of the insulator. Direct field emission into the conduction band of trapped and valence band electrons also contributes to the enhanced emission, but for simplicity we consider only field emission from the metal.

In view of the discussion in Sec. II, the magnitude of the electric field through the insulator is assumed to have the form shown in Fig. 1. Here, space charge in the dielectric interior, except near the vacuum surface, is neglected. At the other extreme, the principal field variation may occur near the dielectric/metal interface, in which case the vacuum field extends through most of the dielectric volume. Whether either extreme or some intermediate case prevails depends on the details of the physical processes occurring within the dielectric.

The rate of buildup of internal field is given (assuming a high mobility for electrons and a low one for ions) by

$$\epsilon_o \frac{d}{dt}(KE_i - E_o) = j_i + j_{FN}(e^{ad}p - 1) \tag{6}$$

with an emission current

$$j_e = j_{FN} e^{ad} p \tag{7}$$

where[8]

$$J_{FN} = 1.5 \times 10^6 \times 10^{4.52\phi^{-1/2}} \times E_i^2 \times 10^{-30\phi^{3/2}/E_i} \mathrm{amp/cm^2} \tag{8}$$

is the Fowler-Nordheim emission into the conduction band lying ϕeV above the Fermi level. In the above equations, E_i and

E_o are in MV/cm, and

$$\epsilon_o = 8.85 \times 10^{-14} F/\text{cm} \times 10^6.$$

The quantity $e^{\alpha d}p$ allows the possibility of positive feedback, leading to continuously increasing emission currents and eventually to the striking of an arc; $e^{\alpha d}$ is the number of conduction electrons produced by collisional ionization as an electron emitted from the metal crosses the thickness d of the insulator, and p is the probability that an electron arriving at d is emitted into vacuum. The ionization coefficient αE_i is a function of the internal electric field E_i and $p(E_o)$ a function of the vacuum electric field. To gain some notion of magnitudes, we adopt a result given in O'Dwyer for $A\ell_2O_3$[9]

$$\alpha = 4 \times 10^7 \exp(-14/E_i)\text{cm}^{-1} \tag{9}$$

At the level of a few MV/cm, substantial ionization may occur in a layer of $d \sim 1000$ Å. Specific numbers on the emission probability are not known, but to be explicit, and in view of the discussion in Sec. III, we shall take it to be of the form

$$p(E_o) = e^{-E_c^2/E_o^2} \tag{10}$$

with E_c perhaps of order 0.1 – 1.0 MV/cm. Equation (6) contains features in qualitative accord with observed results, namely 1) a voltage threshold, 2) prebreakdown emission currents, and 3) an arc rate.

IV. Voltage Threshold

Consider the currents in Eq. (6)

$$j_1 = j_i + j_e \tag{11}$$

$$j_2 = j_{FN} \tag{12}$$

which we show graphically in Fig. 2 as a function of E_i for fixed values of E_o. In Fig. 2a, the emission is small, and perhaps $\alpha d \ll 1$ as well. In this case, the intersection at A is a stable equilibrium in which the incident ion current and any emission are balanced by j_{FN}. In Fig. 2c, no equilibrium is possible. In identifying E_{ot} (Fig. 2b) with a voltage threshold, we are tacitly assuming that the voltage drop across the insulator is small so that E_{ot} will be related to the difference between applied bias and coverslip potential. Threshold values of E_o and E_i can be calculated from the conditions

$$j_1(E_{it}, E_{ot}) = j_2(E_{it}, E_{ot})$$

$$\frac{\partial j_i}{\partial E_i}(E_{it}, E_{ot}) = \frac{\partial j_2}{\partial E_i}(E_{it}, E_{ot}) \tag{13a}$$

Observe in addition that for a fixed value of E_o, we may also obtain a threshold value for the ion current. If we take p from Eq. (10) and set

$$\alpha = \alpha_o \exp(-H/E)$$

$$J_{FN} = AE^2 \exp(-B/E)$$

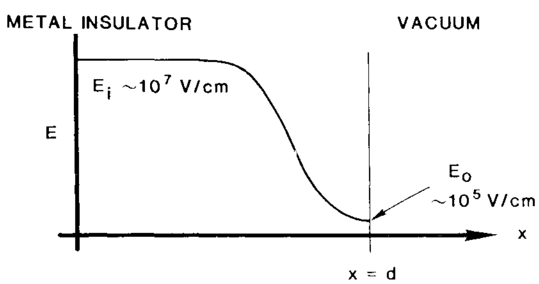

Fig. 1 Electric field profile in insulator.

then Eq. (13a) leads to the following relations involving threshold fields and ion current

$$j_i = j_{FN}\left(1 + \frac{2x+1}{(H/B)\alpha_o d}\exp(H/Bx)\right)^{-1} \tag{13b}$$

$$\left(\frac{E_o}{E_c}\right)^2 = \alpha_o d \exp(-H/Bx)$$

$$+ \ell n[1 + (H/B\alpha_o d/2x+1)\exp(-H/Bx)]/1 \tag{13c}$$

where, for convenience, we have omitted the subscript t denoting threshold values.

V. Prebreakdown Emission Current

Synder[1] observed in his experiments that as a precursor to arcing the coverslips swing a few hundred volts negative in a few hundred seconds in a plasma of density $n \sim 10^4 \text{cm}^{-3}$. In order for this to occur, the emission current collected on the coverglass must exceed the ion current. Moreover, to lead to arcing, we require $e^{\alpha d}p > 1$. To be consistent with a slow charging rate, however, it is required that the rate of buildup of E_i during the prebreakdown phase be slow, with a time scale on the order of hundreds of seconds, about the same as that of the original positive charging of the coverslip by ions. Moreover, from Snyder (Ref. 1, Fig. 5a), it appears that nearly all the cell coverslips are going negative at once, leading to the conclusion that the total emission current is only a few times the ion current. With an interconnect fraction area of 0.05 and the assumption that a significant fraction of the emission goes to the coverslips, it appears that the emission current density at an interconnect is also only a few times the incident ion current at that location; this follows if focusing increases the ion current density at the interconnects by an order of magnitude or so above the areal average. Taking $e^{\alpha d}p$ to be a few times unity, $\phi = 2eV$ and $j_i = 10^{-9}$ amp/cm^2 gives an internal field of

$$E_i \sim 4.6 \text{ MV/cm}$$

The time scale for E_i field buildup at this current level is

$$t_E = E_i/(dE_i/dt) \tag{14}$$

while the time scale for emission current buildup is somewhat shorter,

$$t_p = j/(dj/dt) \tag{15}$$

Specific numerical values are arguable, but the important point is that because of the exponential dependence of currents on fields, emission currents can exceed ion currents for substantial periods of time before they begin to increase at a rapid rate, say on a microsecond time scale. In fact, in circumstances corresponding to Fig. 2a, a steady emission current

$$j_e = (j_i e^{\alpha d}p)/(1 - e^{\alpha d}p) \tag{16}$$

in excess of j_i is possible.

VI. Arc Rate

We see from Eq. (8) that only slight increases in E_i are required to bring emission currents into the ampere range. In the example of the previous section, increasing E_i (4.6 MV/cm) by a factor of two increases the emission current by eight orders of magnitude, with a corresponding reduction in the time scales for internal field and current growth. Accordingly, we may define the interval between arcs, or the inverse arc rate

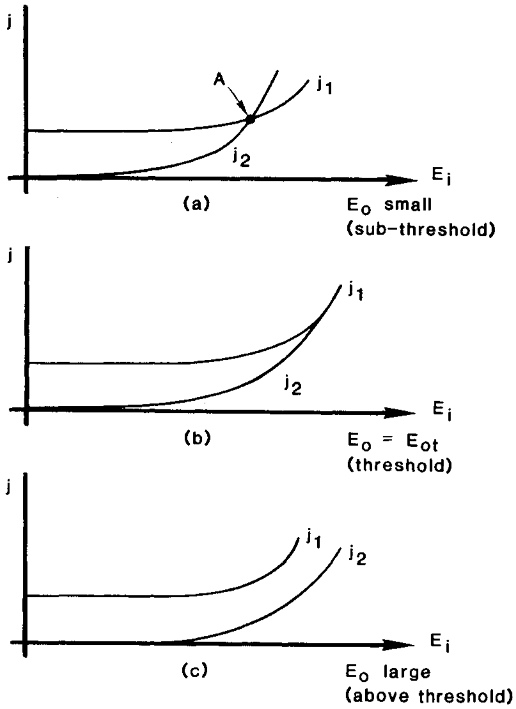

Fig. 2 The stable, unstable, and threshold conditions.

R^{-1}, as

$$R^{-1}=T=\int\frac{K\epsilon_o dE}{j}\approx K\epsilon_o\int_{E_i(o)}^{E_{i\max}}\frac{dE}{j_i+j_{FN}(e^{\alpha d}p-1)}\tag{17}$$

assuming that we are in the situation corresponding to Fig. 2c. Here, $E_{i\max}$ is a field large enough that growth time scales are much shorter than intervals between arcs; for all practical purposes, we may set $E_{i\max}=\infty$. The quantity $E_i(o)$ is the "initial" value of internal field, or in the case of repetitive arcs, the value to which E_i relaxes after the completion of a discharge. Observe that R is an increasing function of j_i for fixed E_o.

To be precise, the term dE_o/dt in Eq. (6) should not be ignored. The arc rate defined by Eq. (17) is strictly appropriate only for fixed external fields E_o and does not take account of the charging of surfaces by emission currents.

The variation of T with E_o in the vicinity of E_{ot} can be estimated. For this purpose, we consider variations that are small in the sense that

$$\left|\frac{E_o-E_{ot}}{E_o}\right|\ll 1\tag{18}$$

but large in the sense that

$$m=\frac{p(E_o)}{p(E_{ot})}\gg 1\tag{19}$$

This is possible if in Eq. (10) $E_c\geq E_{ot}$. Now express T in the form

$$T=K\epsilon_o\int_{E_i(O)}^{\infty}\frac{dE_i}{j_i+j_{FN}(e^{\alpha d}mp_t-1)}\tag{20}$$

$$=\frac{K\epsilon_o}{m}\int_{E_i(O)}^{\infty}\frac{dE_i}{j_{FN}e^{\alpha d}p_t+j_i-j_{FN}/m}\tag{21}$$

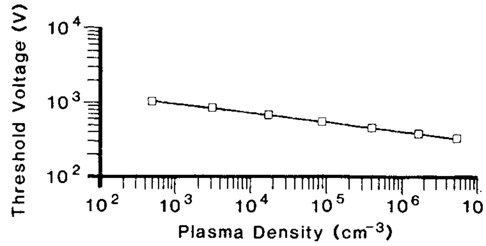

Fig. 3 Threshold external electric field vs plasma density (theoretical curve).

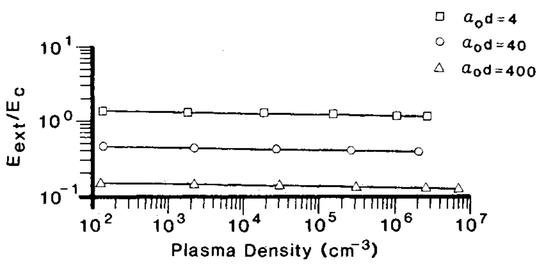

Fig. 4 Threshold voltage vs plasma density ($E_c=700$ V/cm, $\ell=.01$ cm) $\alpha_o d=40$, $H=14$ MV/cm, $B=65$ MV/cm.

(We note that for $m\rightarrow l$, $T\rightarrow\infty$, is a property of the threshold curves of Fig. 2b and not related to the ∞ upper limit of the integral.) If $E_i(O)$ is large enough that the first term in the denominator of the integrand dominates the integral, then

$$R=\frac{p(E_o)}{K\epsilon_o}\left(\int_{E_i(O)}^{\infty}\frac{dE_i}{j_{FN}e^{\alpha d}}\right)^{-1}=\frac{e^{-E_c^2/E_o^2}}{K\epsilon_o}A\tag{22}$$

where A is independent of E_o if $E_i(o)$ is also. Thus, for $E_c>E_o$, R is a rapidly rising function of the vacuum field at the insulator surface. If $E_i(o)$ is small, R will be dominated by the time required to build up to near threshold conditions:

$$R\simeq j_{ion}/K\epsilon_o E_{ot}\tag{23}$$

Between those two extreme cases, the dependence on j_i and E_o will be more complicated.

VII. Numerical Results

So far results have been obtained only for the threshold values of internal and external electric fields. Calculations of arc rate are of interest but have not been brought to a satisfactory state because of uncertainties in the values to which internal and external electric fields relax following a discharge pulse.

Plots of E_{ot}/E_c vs density found by solving Eqs. (13a) and (13b) are shown in Figs. 3 and 4 for a few values of the parameters H/B and $\alpha_o d$. The plotted curves are significant in that for the considered parameters, the external threshold fields (or voltages) are slowly varying functions of plasma density, in accord with the experimental curve shown in Fig. 5.[10] To determine a voltage value from these curves, we set $E_{ot}=V_{ot}/\ell$ where $\ell\simeq.01$ cm is the distance scale on which the potential around the interconnect drops by an appreciable value. With $E_{ot}/E_c\simeq0.4$, as in Fig. 3, the value of E_c can be deduced to be

$$E_c\sim\frac{400}{0.4\times0.01}=10^5\text{V/cm}$$

The curve in Fig. 4 was obtained with $\phi=0.7$ eV and $H=14$ MV/cm such that for $n=10^4$ cm^{-3}, $E_t\ell=V_t=700$ V. These values permit one to fit Stevens' data quite well but require a value E_c of only 700 V/cm.

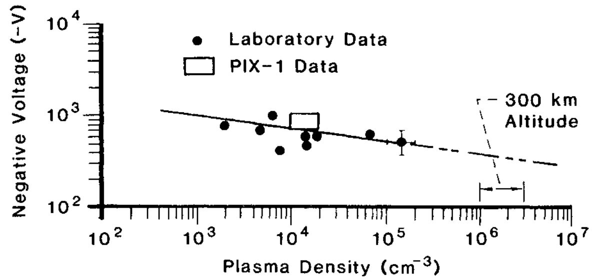

Fig. 5 Voltage threshold for breakdown (from Ref. 10).

If Eq. (22) is valid with an initial value $E_i(o)$ independent of external field, the implication is that the arc rate varies strongly with voltage only for E_{ot}/E_c less than unity. It is conceivable, however, that $E_{ot}/E_c>1$, in which case the strong dependence of the arc rate on the voltage would require the initial internal field $E_i(o)$ to depend on the value of the external field that exists following a discharge pulse.

VIII. Summary and Conclusions

We have attempted to develop a theoretical basis that is in substantial accord with the observed characteristics of arc initiation summarized in Sec. I. The theoretical basis consists of the following elements: 1) an insulating layer over metallic interconnects, 2) the charging of a thin but finite region of the layer by incident ions, 3) the injection of electrons into the insulator's conduction band by Fowler-Nordheim emission from the metal, 4) heating and further ionization by the interior electric field, and 5) the thermionic emission of conduction band electrons into vacuum with an effective temperature related to the surface electric field. Qualitatively, the proposed theory enables the determination of threshold voltages conforming to observations; quantitatively, the relevant properties of insulators are hard to determine, especially when the insulator itself is unknown. The implications of the theory are of practical interest, however, even in the absence of such information.

For example, the theory predicts that arcing is related to the magnitude of the vacuum field at the insulator surface. Arcing should depend, therefore, not only on the applied bias, but also on the geometry of the interconnect. It is conceivable then that the geometry could be effected to reduce the arc rate or, optimistically, eliminate arcing altogether.

For repetitive arcing to occur, as envisioned in Sec. VI, it is necessary that successive discharges occur at different locations on the array, or if this is not the case, that the contaminant layer either survive the discharge or be restored, by deposition, say, in the interval between discharges. The former alternative is almost certainly the prevailing one. It is nonetheless of interest to inquire into the fate of a contaminant layer that has conducted a discharge. In the experiments conducted by Snyder[1,2] and Stevens[10] it is probable that the layer survives the discharge. This assertion is based on the estimate $j_1E_i\Delta t$ of the energy density in the layer following a discharge of duration Δt. Assuming $\Delta t\sim10^{-6}$ s, a discharge current of 0.1 A[1,2] over an area of 10^{-3} cm^2 and a field of 10^6 v/cm

$$j_1E_i\Delta t\sim10^3\,\mathrm{J/cm^3}$$

which is sufficient to raise substantially the temperature of the layer but probably not sufficient to vaporize it. Moreover, it is reasonable to suppose that the deposited energy is transported away from the layer in the rather long time interval before the layer conducted a second discharge current.

It is also clear, however, that currents and pulse durations not very much greater than those observed by Snyder[1,2] and Stevens[10] ($j_1E_i\Delta t\sim10^5$ J/cm^3) would be sufficient to vaporize and ionize the contaminant layer, thus forming an ionized gas that would expand into vacuum. At this point the nature of the discharge would differ dramatically from that considered in this paper. In fact, as the capacitance in the discharge circuit increases, the discharge duration Δt and the peak discharge current both increase. Miller[3] conducted experiments in which the power supply and its capacitance were not isolated by a limiting resistor from the discharge circuit. The currents (>1 A) and pulse durations ($\Delta t\sim10^{-3}$ s) were accompanied by clear evidence of material damage in the region of the metallic interconnects.

So far we have addressed only the question of the relatively slow process of arc initiation which, under the experimental conditions referred to here, occurs on a time scale ranging from seconds to hundreds of seconds. We expect to consider the rapid process of arc buildup, which occurs on time scales in the microsecond or shorter range, as well as questions of arc limiting in future work.

Acknowledgment

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References

[1]Snyder, D. B., "Discharges on a Negatively Biased Solar Array," NASA TM-83644, Oct. 1983.

[2]Snyder, D. B., "Characteristics of Arc Currents on a Negatively Biased Solar Array," NASA TM-83728, June 1984.

[3]Miller, W. L., "An Investigation of Arc Discharging on Negatively Biased Dielectric-Conductor Samples in a Plasma," NASA TM-83787, Oct. 1983.

[4]Jongeward, G. A., Katz, I., Mandell, M. J., and Parks, D. E., "The Role of Unneutralized Surface Ions in Negative Potential Arcing," presented at the 1985 IEEE 22nd Annual Conference, NSRE, Monterey, CA, July 22–24, 1985.

[5]Latham, R. V., "Prebreakdown Electron Emission," *IEEE Transactions on Electrical Insulation*, Vol. EI-18, June 1983.

[6]Rose, A., "Recombination Processes in Insulators and Semiconductors," Physics Review, 1954.

[7]Lambert, M. A. and Mark, P., *Current Injection in Solids*, Academic Press, New York, 1970.

[8]Latham, R. V., *High Voltage Vacuum Insulation: The Physical Basis*, Academic Press, London, England, 1981.

[9]O'Dwyer, J. J., "The Theory of Electrical Conduction and Breakdown in Solid Dielectrics," Clarendon Press, Oxford, England, 1973.

[10]Stevens, N. J., "Interactions Between Large Space Power Systems and Low-Earth-Orbit Plasmas," Spacecraft Environmental Interactions Technology Conference, Colorado Springs, CO, NASA CP-2359, AFGL-TR-85-0018, 1983, p. 263.