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THEORETICAL MODEL FOR PLASMA OPENING SWITCH

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Abstract

The theory of an explosive plasma switch is developed and compared with the experimental results of Pavlovskii and work at Sandia. A simple analytic model is developed, which predicts that such switches may achieve opening times of approximately 100 ns. When the switching time is limited by channel mixing it scales as $t = C(m d_0)^{1/2} P_0^2 P_e^{-5/2}$ where m is the foil mass per unit area, d_0 the channel thickness and P_0 the channel pressure (at explosive "breakout"), P_e the explosive pressure, C a constant of order 10 for c.g.s. units. Thus faster switching times may be achieved by minimizing foil mass and channel pressure, or increasing explosive product pressure, with the scaling exponents as shown suggesting that changes in pressures would be more effective.

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I. Introduction

Fast opening switches of the type described by Pavlovskii et al.¹ are of interest for pulsed power applications. Such switches are needed, for example², with inductive generators which must have a shorting load while they build up a magnetic field. At peak flux this short should be switched off and power delivered to the useful load (see e.g., Fig. 2). In this way the useful load does not have to possess a low impedance or be subjected to a slow current buildup for many microseconds. These devices can switch megamperes and megajoules in fractions of a microsecond. Results from computer simulations and experiment show that simple one-dimensional compression of the conducting channel is insufficient to account for the experimentally observed increase in resistance. Consequently, we argue that the Rayleigh-Taylor instability plays an important role in the operation of the device, mixing cool material into the hot conducting gases of the plasma channel, and thus greatly increases its resistivity. A simple analytic model for the behavior of the one-dimensional compression and the growth of instabilities and mixing is developed which gives design criteria that optimize switch performance. Experimental verification is hoped for in the next series of experiments, which will test the scaling laws predicted by the theory developed here.

II. Behavior and Mechanism of Switch

Our understanding of the switch is gained from comparison of experiment and numerical simulation.

The basic geometry of the Pavlovskii switch¹ is shown in Fig. 1. Initially a thin Mylar (Dacron) foil has its outer surface coated with a thin aluminum layer. The large current rapidly forms a conducting plasma channel, either by striking an arc in the air of the channel, by vaporizing the aluminium, or both. We for the moment assume that in the Russian work the channel was initially filled with air at STP. This seems reasonable, since the granular explosive used could not support a vacuum, and the thin (no dimensions given) Mylar foil could not support a substantial pressure differential. It would seem likely, based upon experiments done at Sandia by D. McDaniel and R. Stinnett³ as well as computational work by L. Baker,⁴ that an air discharge would be initiated. For our purposes it does not matter, since the mass of air greatly exceeds that of the aluminium vapor. We will discuss below possible effects of diminished gas pressure and/or the use of different fill gases in the switch. For our purposes we will model the device by assuming that initially we have a hot plasma channel of air that was initially at atmospheric density and pressure.

G. Sauve⁵ has integrated the experimental data in Ref. 1 on the derivative of current to give an approximate $I(t)$ curve. This may be used to estimate the energy deposited in the device, and hence the state of the plasma channel. The results are shown in Table I. We have assumed¹ constant resistance of 0.8 milliohms in obtaining the energy density, and dimensions are from an analysis of G. Sauve who finds from Fig. 1 of Ref. 1 a device length of 12 cm and radius of 10 cm. The 100 Angstrom

thick layer of aluminium gives rise to a density of 1.7×10^{-4} g/cc, assuming a radius of 10 cm. The separation* energy is 1.2×10^{11} erg/g, therefore, the aluminum would be fully vaporized at $2. \times 10^7$ erg/cm, that

Table I. Estimates of "action integral" and energy deposited in the Pavlovskii switch.

<u>time</u> <u>microsec</u>	<u>current</u> <u>megamps</u>	<u>action</u> <u>amp²sec</u>	<u>energy deposited per unit length</u> <u>ergs/cm</u>
0	0.	0.	0.
40	1.5	$3. \times 10^7$	2.0×10^{10}
80	2.	1.4×10^8	9.3×10^{10}
100	3.	2.5×10^8	1.7×10^{11}
115	7.	5.5×10^8	3.7×10^{11}

is, instantaneously on the timescale of the experiment, assuming all the energy were deposited in the aluminum. The Mylar has a density of 1.4 gm/cc and a separation energy of 2×10^{11} erg/g. Assuming that the Mylar is roughly 1 mil thick, this would give a density per unit length of roughly .22 gm/cm, requiring $4. \times 10^{10}$ erg/cm or roughly 75 μ s.

to vaporize the Mylar. It therefore seems likely that the Mylar will be more or less fully vaporized by the time of switching action, but for the bulk of the experiment the mylar is not part of the conducting channel. For a gap width of 2.5 mm, the aluminum contributes a mass density of 3.74×10^{-6} gm/cc in the channel, that is, about 1/300 of the ambient air; the ratio of particle number density for air and aluminum is then also roughly 300 to 1, i.e., the aluminium atoms are a slight impurity if mixed with the air. At later times the density may rise due to

*The separation energy is the energy required to separate the atoms to infinity, i.e., it is the sum of the heats of fusion and vaporization and the integral of specific heat multiplied by dT.

material ablated off the outer wall and the Mylar. The numerical simulations suggest that the amount of ablated material is not large, but this result is not considered conclusive due to approximations made in the computer code. These calculations are Lagrangian, and consequently cannot simulate the mixing of such material into the plasma. There is not enough time to allow such mixing, and by the time the Mylar is vaporized, the air channel is a sufficiently good conductor that little current, if any, will be shunted through the much more resistive Mylar vapor. It is reasonable, therefore, that such material does not contribute significantly to plasma density within the channel or provide a significantly enlarged or denser conducting region.

Numerical simulations were performed with a version of the CHARTB⁶ magnetohydrodynamics code. This code is capable of treating the detonation of the explosive self-consistently, using the J-W-L equation of state⁷ for the material, the phase transitions of the Mylar, the ablation of the Mylar through heat conduction, the magnetohydrodynamics of the plasma, and the external circuit. These calculations were performed using data for the explosive PETN. The calculations tend to be rather costly. One must resolve the plasma channel, requiring about a dozen zones within the air/aluminium vapor region. Mass matching to the Mylar layer results in extremely fine zones in the Mylar, resulting in restrictive timestep limits due to the Courant-Friedrichs-Lewy criterion (i.e., that a sound wave should not travel more than one zone per time step) in that material. Differing models were used for the external load. Since the existence of such a load would provide another current path, and hence (during and after switching) diminish the Joule heating, omitting the load in the

simulations might cause pessimistic switching behavior. It is dangerously inconsistent to impose a current through the system, since this could result in excessive Joule heating and therefore quite different device behavior. Therefore the current source was modeled as a stray inductance of 38 nH in series with a generator inductance of the form $60\text{nH}(1-(t-t_0)^2/T^2)$, where $t_0=10$ microsec, and $T=20$ microsec.(Fig.2). For the initial condition, $I=4$ megamps is assumed. For the load, we assume an inductance of 30 nH which is switched in across the switch as soon as a voltage drop of 100 kV is exceeded. These parameters are based on data in Ref. 1 and typical behavior of magneto-explosive generators.^{2,8} For initial conditions in the plasma channel, a temperature of 2.5 eV of the air-aluminum channel is assumed. Results are shown in Figs. 3 and 4.

Such one-dimensional simulations show much poorer switching action than observed. The channel expands somewhat, pushing the Mylar layer inward; the Mylar has not significantly melted or vaporized by the time the explosive front has reached it. The code results show that predetonation of the explosive by impact of Mylar (either solid or vaporized) is highly unlikely in plausible designs. The explosive soon compresses the channel by a factor of about twenty from its most expanded state (initially, .25 cm in radius, expanded by about a factor of 3 to .85 cm radial extent at about 7 microseconds, and compressed to minimum radius of .04 cm at 10.1 microseconds). The resistivity of the channel is reasonably constant throughout the current rise and a good match to the experimental value of 0.8 milliohms; it starts at 2.5 milliohms and drops to 0.8 milliohms at 7 microseconds. At its highest, during switching, it achieves only 3.4 milliohms, or a factor of about 4, compared to the observed experimental value of about a factor of 200 (peak resistance

200 milliohms). A plot of the gap size d , channel resistance, and current are shown in Fig. 2. This simulation of the Pavlovskii switch will be referred to as the "canonical" case.

There are a number of possible explanations for this discrepancy, including defects in the numerical simulation. Spitzer resistivity (including a neutral collision term) is used, and while this may be reasonably accurate during most of the simulation, in the higher density compressed phase, the plasma becomes "non-ideal" (there are few particles per Debye sphere), resulting in an underestimate of the resistivity. We doubt that such effects can account for a factor of 100 increase in resistivity, although this possibility cannot be completely ruled out due to our ignorance of the proper expressions for resistivity. D. L. Boercker⁹ is currently working on such effects, and we hope to be able to apply his resistivity models to the problem.

One can readily show that enhanced thermal losses from the thinner channel cannot account for the rapid resistance rise, assuming classical thermal conduction and loss mechanisms. Wall ablation is not simulated in these calculations, but results of other calculations and the Mylar ablation rate suggest it is unimportant. Due to uncertainties in the physics, the modeling of the conduction coefficient, loss mechanisms, and ablation may be a source of error which cannot entirely be ruled out.

We consider a more plausible explanation to be that the actual device operation is not one-dimensional. It is clear that, near the "turnaround" of the plasma channel implosion by the Mylar-explosive product gas combination, we have a Rayleigh-Taylor unstable situation, i.e., the classic case of the light (channel gas) material decelerating the denser material. This would cause "spikes" of dense cool material

to penetrate into and mix with the channel material. The materials would then mix, hot channel material would be cooled and the resistivity would rise. Note that it is almost certain that such instabilities must develop if the mylar at this time is denser than the air channel gas. The only question is whether such mixing is important in switch operation. Since the one-dimensional model does not account for the resistance rise, this mechanism is a very plausible means of doing so. The Mylar mass, as noted above, is 50 times that of the channel, and code calculations give temperatures of the order of 6 eV in the channel and on the order .1 eV in the Mylar. Under these conditions, mixing on the atomic scale would result in temperatures on the order of .2 eV (the specific heats are roughly equal) and the Spitzer resistivity would rise by a factor of 100 or more, easily accounting for the observations. We have performed additional CHARTB simulations starting near turnaround but assuming a mixed channel. The initial momentum is taken from the unmixed runs, but the channel is mixed and the temperature is reduced to .6 eV (lower temperatures would give an insulated channel given the resistivity model used). Temperature, pressure, and velocity profiles of the mixed channel are shown in Fig. 5. The fully mixed channel has an initial resistivity of 60 ohms (compared to the observed peak of .2), suggesting that only a partial mixing is necessary to account for the observations. The resistivity of the mixed channel drops very rapidly thanks to Joule heating, which raises the temperature to 2 eV. Expansion of the mixed channel brings the resistivity down in the few milliohm regime on a time scale of a few hundred nanoseconds.

The Taylor instability hypothesis for switch operation is plausible. The Mylar reaches a peak velocity in compressing the channel of $2. \times 10^5$ cm/sec,

and is decelerated over one microsecond to rest, giving an average acceleration of 2×10^{11} cm/sec². This deceleration causes the interface between the hot, tenuous channel and "pusher" (Mylar and explosive products) to become Rayleigh-Taylor unstable. The most effective mode for mixing will have a wavelength on the order of the channel thickness at turnaround; we will use 1 mm to allow for the fact that the channel is diminishing in width. This gives a growth rate of 6×10^6 sec⁻¹. If we assume an initial perturbation at implosion time of 10% of the wavelength, i.e., .1 mm, we have the penetration time for .04 cm of 200 ns. This is in good agreement with the switching times of the Pavlovskii device and similar test devices at Sandia. Of course, this is merely an estimate, and a more accurate treatment should consider such complications as:

- 1) The instability starts to grow as soon as deceleration sets in, before turn around; we will show this is a relatively small effect, but tends to cause an overestimate of the switching time.
- 2) The perturbation does not grow linearly as assumed, but will saturate in general and grow more slowly, with approximately constant velocity in our case. This would cause us to underestimate the switching time, but for wavelengths on the order of the required penetration distance, the effect is not important.
- 3) We have not explicitly treated the transition of the instability to turbulent mixing. This would probably result in our underestimating the time for switching action to take place. This would be almost impossible to treat accurately, but we will discuss this at somewhat greater length below. We will also discuss in more detail the assumptions about the character of the initial perturbation.

Experiments done at Sandia using a planar geometry seem to confirm the model for device operation discussed here. The resistance of the switch is observed to rise to an intermediate value, remains rather constant for a few microseconds, and then rises on a similar timescale to peak resistance. This is consistent with the model, the initial small rise in resistivity due to the decrease in channel cross section, and the later rise due to the mixing and cooling of the channel. Figure 6 shows an experimental resistance plot and voltage plots which illustrate this behavior. Tests in the cylindrical geometry of Pavlovskii should provide better switching action and provide more convincing confirmation of the model.

III. Simple Model for Switch Action

A. Preliminary Remarks

A simple model for switch operation (channel compression and mixing) would be of value in providing analytic formulae for switch behavior as a function of the various parameters, enabling optimal parameters to be estimated without costly numerical simulations.

The Lorentz force need not be modeled in detail in such a theory. The geometry of the switch is an "inverse pinch." As in the usual pinch, the Lorentz force tends to compress the conducting channel, but in this case the compressional force squeezes the channel out against the wall. For a perfect conductor, this force would be the usual magnetic pressure force. As some of the field (current) has diffused in, the force is less effective in compressing the plasma. The Lorentz force, therefore, acts in concert, and not against, the explosive pressure, and a portion of the magnetic pressure (not more than half of $B^2/8\pi$) should be added to the explosive pressure. This is generally a negligible term when field penetration is accounted for, since the field has a number of microseconds in which to penetrate the plasma.

The use of this inverse pinch configuration is, we believe, motivated by the need for suppression of MHD instabilities. The device must carry a large current for 100 μ s. As growth times for MHD instabilities are typically a few microseconds, such instabilities could disrupt the channel and consequently interrupt the current if they were present. Therefore, a gas-embedded Z-pinch configuration, which might seem attractive at first, would not work as it would have severe problems with kink and sausage mode MHD instabilities on the time scale of current rise through the device. In the inverse pinch configuration, the plasma channel

is pushed outward, against the insulating wall, and is stable. Therefore, we must use the inverse pinch. This is not a great sacrifice, though one might at first think the compression of a thin gap at large radius would be less effective in rapidly reducing channel cross section than compression of a channel on the axis, because (1) the pressure in the channel would rise more rapidly, with the "effective gamma" (see below) twice as large as in the one-dimensional case, and (2) the switching is not due to a one-dimensional change in cross-sectional area.

The proper modeling of both the detonation and electrical properties of detonation products of the explosives must be considered. One factor that may limit switch performance is the diffusion rate of the magnetic field through the conducting explosive product gas. If this gas were very conducting, it would freeze the current into the plasma channel and therefore prevent switching action. The available data demonstrate that measured values are a function of electrode sheaths, explosive product gas expansion, composition, etc. in a manner that makes for an obscure picture. Pavlovskii et al¹ refer to the work of Brish et al¹⁰; more data can be found in Cook¹¹, and our bibliography contains in addition to cited works a list of papers on the electrical properties of explosive products. Table II compares the average conductivities reported by each. We find that in addition to discrepancies in the magnitudes, Cook finds TNT more resistive than PETN or Tetryl by a factor of about two, while the Russian work finds it less resistive by very large factors. The densities involved were somewhat different; all the Russian work used powdered explosives of density .8; the work of Cook used TNT of density .9, PETN density .91, and Tetryl of density 1.0 gm/cc. Still, the differences

are hard to explain, since results of Cook for TNT of densities 1.03 and 1.57 differed relatively little from the value at .9 gm/cc .

Table II. Comparison of Measured Explosive Product Resistivities

<u>Explosive</u>	<u>Resistivity (Ref.10)</u>	<u>Resistivity (Ref. 11)</u>
PETN	8.	1.13
Tetryl	6.	1.10
TNT	.25	2.03
Comp. B	1.	2.56

All results are in ohm-cm. Ref. 9 averages over first cm of explosive product zone.

The conductive zones behind the explosive front measure about 2-3 cm typically.¹² Using the less resistive and therefore more pessimistic data of Ref. 11, we can estimate the diffusion time, $12 d^2/r$ (ns), where d is the width of the conducting zone in cm and r is the resistivity in ohm-cm. Since the conductivity falls exponentially with distance from the front, we will again be pessimistic and overestimate the required diffusion time by using 1 ohm-cm for the whole front region. The resulting diffusion times (and therefore minimum switching times) are on the order of 50-100 ns.

There is some support for the larger resistivities in other Russian work,¹² which would decrease diffusion times. However, other work^{13,14} indicates that electrical breakdown of the explosive product gas can be important. This too is a function of expansion of explosive products, etc. Ref. 12 finds strengths of order 100-120 kV/cm for PETN, for example, while Ref. 13 finds values on the order of 140 kV/cm and Ref. 14 finds about 600 kV/cm. While the role of electrode sheaths, etc. are discussed

it seems safe to conclude that a maximum of about 100 kV/cm may be sustained before an electrical breakdown through the explosive gases compromises the switch. This imposes another limitation, that the voltage inductively induced upon switching, $L_{\text{ext}} \dot{I}$, must be sufficiently small. In Pavlovskii's work $L_{\text{ext}} = 30 \text{ nH}$, in a device which can stand off about 1 MV across the distance $h = 12 \text{ cm}$. If I is about 10 MA, this means the switching time must exceed 300 ns. In the experiment voltage gradients never exceeded about 25 kV/cm in the switch, so there was no arcing. This constraint may be the most severe on switching time in some applications, and limits the external load inductances for which such switches are useful. Current applications contemplate loads with inductances perhaps a factor of ten lower than Pavlovskii's, and so this constraint will not severely impact them. A limit to $L_{\text{ext}} \dot{I}$ will limit the power $(L_{\text{ext}} \dot{I})I$ which may be delivered; $L_{\text{ext}} \dot{I}$ must be less than $E_{\text{crit}} h = k 0^5 h$. For given design parameters, this fixes a minimum length for the smallest dimension of the device. For example, to deliver 10 MJ of energy in 100 ns with 10 MA of current, the voltage must be 20 MV, implying a device of length and radius of at least 200 cm to avoid a shunting arc.

We also note, while on the subject of explosives, that the explosive pressure falls very rapidly with increase in volume. The J-W-L equation of state⁷ gives the pressure $p = A \exp(-R_1 V) + B \exp(-R_2 V) + C V^{-w}$ where V is the ratio of specific volume to that of the predetonation explosive. For $V=1$, the pressures are generally about 1/3 or so of the Chapman-Jouguet pressure. Since C is very much less than A and B and since the constants R_1 and R_2 are large, only a slight expansion reduces the explosive pressure considerably. Therefore, the explosive pressure used in our

model, and taken from the CHARTB simulations, will be much lower than the Chapman-Jouguet pressure. However, through the inertia of the Mylar "pusher," we will achieve pressures in the plasma at turnaround that exceed the pressure on the other side of the Mylar.

B. Simple Model: Description

We write a simple equation of motion for the Mylar "pusher," treating it as an inert mass slug. The CHARTB simulations show that shocks due to the explosive wave are not important in determining the dynamics or the state of the Mylar or the plasma channel, due to the relatively low pressure of the explosive products by the time they reach the Mylar. Letting d = channel width (cm), d_0 = channel width at time of contact with explosive products, P_e = (average) pressure of explosive products on the mylar, m foil mass/unit area, we write the equation of motion as

$$m \ddot{d} = P_0 (d_0/d)^g - P_e \quad (1)$$

where P_0 is the gas pressure in the conducting channel (assumed uniform) at the time of explosive product contact with the Mylar, and g is an "effective gamma" for the compression of the gas. If the channel were an ideal, ionized gas, g would be about 5/3. Since the compression increases the current density for constant current, the additional Joule heating tends to increase g above this value. On the other hand, any rise in temperature will result in increased ionization at these temperatures. Also, losses due to processes such as free-free and bound-free emission in the optically thin plasma should result in heat loss that will tend to reduce the value of g . From the CHARTB simulations, which include such loss mechanisms, we find $g=1.5$ is a good approximation.

Using the dimensionless variables $x = d/d_0$ and $T = (P_e/d_0 m)^{1/2} t$ with $r = P_0/P_e$, we have (the raised dot now denoting derivative with respect to T):

$$\ddot{x} = r x^{-g} - 1 \quad (2)$$

For g not equal to 1, this integrates to

$$\dot{x} = -(2r/(g-1) (1-x)^{1-g} + 2(1-x))^{1/2} \quad (3)$$

where we have used the initial condition $x=1, \dot{x}=0$ (at $T=0$).

One can prove that turn-around always occurs before $x=0$.

For an "isothermal" compression $g=1$ and we find

$$\dot{x} = (2r \ln x + 2(1-x))^{1/2} \quad (4)$$

The condition for turn around is then $x_{ta} = 1 + r \ln x_{ta}$. For small r we expect that x_{ta} will be small compared to 1, and therefore $x_{ta} = \exp(-1/r)$.

The cases $g=1.5$ and $g=2$ may be relatively simply solved for x_{ta} . The latter case, which would apply if we were compressing a magnetic field, yields $x_{ta} = q/r$. The value $g=1.5$, which seems appropriate from the ratios of d_0/d_{ta} and P_{ta}/P_0 in the CHARTB simulations, gives $x_{ta} = .5(2q+1-(4q+1)^{1/2}) = .5(4r+1-(8r+1)^{1/2})$. For small r this gives $4r^2$ as a simple approximation, which we will use. A comparison shows this is quite good for r less than about .08:

r	'exact' x_{ta}	$4r^2$
.1	.029	.04
.08	.019	.0256
.06	.012	.014
.01	.00038	.0004

Note it always overestimates x_{ta} ; this will lead to a slight amount of pessimism about performance.

Using the results of the numerical simulation of the "canonical" case, we can estimate g . Assuming a $PV^g = \text{constant}$ law, with V (the specific volume) proportional to d , a fit to d_0 and d_{ta} with P_0 and P_{ta} gives $g=1.5$. Assuming $TV^{g-1} = \text{constant}$, we find $g = 1.3$, so the gas does not behave exactly polytropically. Since we are concerned with dynamics, i.e., the behavior of the pressure, we will use $g = 1.5$ since it may be modeled quite simply, as we have done above. The maximum gas pressure at the Mylar-explosive interface is $3.6 \times 10^9 \text{ dynes/cm}^2$ (3.6kBar), in a fairly linearly rising pulse, suggesting that we should take P_e to be roughly half this, or 1.8×10^9 . Then the parameter $r = 0.10$, using $P_0 = 1.6 \times 10^8$, which yields $x_{ta} = .03$, as compared to $x_{ta} = .05$ computationally. Some of the discrepancy may be explained by the fact that the channel is expanding when hit by the explosive, and this additional energy is not accounted for in our simple model which assumes the velocity is initially zero. This additional internal energy would tend to increase x_{ta} . Therefore the model seems to be reasonable; we are currently investigating the possibility of allowing P_e to be a simple function of time, but this seems to require a numerically integrated model (the "trick" used to integrate (2) to find x will not work if t appears in the equations explicitly), and so the model then loses much of its appeal.

Having developed a reasonable model for the one-dimensional compression of the plasma channel, we now turn to consideration of the growth of Taylor instabilities. Such instabilities probably have their principal source in slight irregularities in the explosion front, since the explosive surface is not perfectly smooth. Furthermore, studies of detonation suggest that it is never a plane front, but involves

"turbulence" or three-dimensional motions. These irregularities provide the "seed" for the irregularities that grow due to Rayleigh-Taylor "spikes" that eventually mix with the channel. One would expect, due to the chaotic character of the initial seed, that all wavelengths will be present. The most "destructive", for our purposes, would be one with a scale somewhat larger than d_{ta} . This is because longer wavelengths would grow more slowly, while short wavelengths would "saturate" and cease growing exponentially before they approached amplitude d_{ta} . This argument in greater detail can be found in the literature¹⁵.

Let us now estimate the instability growth during deceleration to turn around. This will be shown to be relatively small. The total growth during the run in will be, approximately,

$$z_{ta} = z_0 \exp\left(\int (ka)^{1/2} dt\right) \quad (5)$$

where z is the perturbation amplitude, a is the deceleration of the mylar, k the wavenumber of the perturbation, $k = 2\pi/\lambda$ where λ is the wavelength, and the integral is over the range of time that deceleration is occurring. We have assumed linear theory applies and used a WKB analysis since a is not constant. Then

$$(ka)^{1/2} dt = (k_0 d)^{1/2} \int_{4r^2}^{r^{2/3}} \frac{dx (x(rx^{-1.5}-1))^{1/2}}{(4r(x^{1/2}-1)+2(1-x)x^{1/2})^{1/2}}$$

where we have used the fact (see Eq.(2)) that deceleration begins at $x=r^{2/3}$ for $g=1.5$. Transforming to the variable $x=s^2$, $s=x^{1/2}$,

$$(ka)^{1/2} dt = (k_0 d)^{1/2} \int_{2r}^{r^{1/3}} \frac{2s^{1/2} (r-s^3)^{1/2} ds}{((1-s)(-4r+2s(1+s)))^{1/2}}$$

Note that we require r be greater than $8r^3$ (i.e., $r < .26$), due to our use of the "small r " approximation for x_{ta} . We now approximate by noting that the major contribution comes near turn around; we replace s with $2r$ in $r-s^3$ and rewrite $2s+4s^2-4r$ as $2(2r)+2sr-4r=2sr$. Then we find the integral becomes $2(1-8r^2)^{1/2} ((1-2r)^{1/2} - (1-r^{1/3})^{1/2})$, where the term involving $1-8r^2$ may safely be set to 1. This gives values for the integral as written above of .324 for $r=.1$ and .307 for $r=.05$; since the growth is proportional to the duration of deceleration and only to the square root of acceleration, we see that smaller r gives slightly reduced growth during run in. Numerical integration of (2) compares well with this estimate, being generally slightly larger (about .4). If we take the growth factor as $.3(kd_0)^{1/2}$, so that at turn around z , the instability amplitude, is $\exp(.3(kd_0)^{1/2})$ times its value at $T=0$, we see that for realistic d_0 and k the effects are rather small; $\exp(.75(d_0/\lambda)^{1/2})$ will be roughly 3 - 4 but not orders of magnitude. Hence the major instability growth and mixing will begin with the "coasting" of the spikes at turnaround. This is fortunate, not a cause for despair, since if the instability grew during run in to a large degree, we would have a great deal of mixing resulting in 1) larger d_{ta} and 2) a more gradual rise in resistivity.

As the instability starts out as a velocity perturbation, it will have a time dependence during linear growth of the form

$$z = z_0 \sinh (Gt) \quad (6)$$

the velocity perturbation being initially $z_0 G = \dot{z}_0$. The error made above by using $\exp$ instead of $\sinh$ is inconsequential. The time for instability growth to size d_{ta} will be $t = d_{ta} / (z_{ta} G_{ta})$ where $G_{ta} = (k d_0 (\ddot{x}_{ta}))^{1/2}$. We have assumed after turn around, the instability spikes coast and mix. For $g=1.5$, we may approximate G as $(k P_e / 8 m r^2)^{1/2}$. The initial amplitude we expect for the turbulent "seed" perturbation as discussed above to be on the order 10% of the wavelength. We have $d_{ta} = d_0 x_{ta} = 4 r^2 d_0$. Thus the growth rate is

$$G_{ta} = (k d_0 (r x_{ta}^{-1.5} - 1) P_e / (d_0 m))^{1/2} \quad (7)$$

and, since 1 is smaller than $1/(8r^2)$, we have $G_{ta} = (\pi)^{1/2} / (4r^2)$. Thus $t = 16 d_0^{3/2} (P_0/P_e)^4 (m/P_e \pi)^{1/2} / z_{ta}$. If we now take z_{ta} to be about 3-4 times z_0 , given the results above, and z_0 to be .1 (e.g., Ref.13), we then have as our approximation:

$$t = 40 / (\pi)^{1/2} d_0^{1/2} (P_0/P_e)^2 (m/P_e)^{1/2} \quad (8).$$

Eq. (8) is our scaling law which illustrates the role of the various parameters on the growth time. Note the desirability of reducing r , that is P_0/P_e , and the weaker dependence upon m , the foil mass per unit area, which also should be minimized. For typical parameters ($P_e = 2 \times 10^9$ (dyne/cm²), $r=.1$ using for P_0 the channel pressure at the time of maximum channel width, that is, the time at which P_e is found and the time at which compression of the plasma channel begins $m = 10^{-3}$ g/cm², $k = 100$ cm⁻¹, $d_0 = 1$ mm) t is roughly 100 ns.

Another way to estimate the mixing time, often used in the analysis of targets for inertial confinement fusion, is the "free-fall" line. This

assumes that the "spikes" of the perturbations move with the fastest velocity attained by the interface (before its deceleration to turn around). In target design one often assumes that the mixing will prevent burn if the free-fall line crosses half the fuel volume before ignition, but let us be conservative and assume mixing will not affect resistivity significantly until the line crosses the entire channel width. Thus we estimate the mixing time as $t = x_m / \dot{x}_m$, where x_m and $\dot{x}_m$ are evaluated for maximum $\dot{x}_m$. Compared to $t = d_{ta} / (z_{ta} G_{ta})$ above we see the answer will probably be fairly similar. By maximizing the magnitude of $\dot{x}_m$ using Eq. (3) we find that $x_m = (3r)^{2/3}$ for $g=3/2$. For $r=.1$, $x_m=.448$ and $\dot{x}_m = -.9519$, giving $t = .47$ nondimensionalized, or using the "typical" values for d_{0,m,P_e} , $t = 100$ ns. This is in good agreement with the results above, as might be expected from the use of a time estimate of the form $t = d_{ta} / (z_{ta} G_{ta})$. This simple model does not tell us, for example, how much the instability grows on the way in (i.e., before turn around); this is of importance in determining the rate of resistance rise.

Conclusion

We have shown that instability growth is rather negligible until turn around; it is the coasting after turnaround that causes mixing. Kelvin-Helmholtz instability of the spikes falling toward the outside wall, as well as the "splattering" as those spikes contact or come very close to the wall, provide the mixing. This development is the hardest to estimate in terms of timescale, since the situation is quite turbulent. The presence of a nearby wall suggests that shear-flow turbulence is appropriate, and therefore that mixing-length theory would apply, and therefore we would have timescales for mixing of order d/v , with d of order d_{ta} and v of the order of the spike velocities. A more detailed discussion of mixing may be found,¹⁶ referring to atomic mixing, estimating t_{mix} as $d^2/(Lv)$ where L is a mean free path, v a velocity (in Ref. 16, a thermal velocity, here a turbulent one about equal to that of the instability). In our turbulent case, $d = L = d_{ta}$, approximately, so $t_{mix} = d_{ta}/v_{rt}$. These estimates would again suggest timescales on the order of 100 ns as above. Since the spikes mix as they "fall" toward the wall, this is still consistent with a time scale of at least 100 ns for switching, perhaps longer, for the parameters discussed.

The switching time may then be estimated as

$$t = \max(t_{diff}, t_{arc}, t_{rt} + t_{mix}) \quad (9)$$

where t_{diff} refers to the time for field diffusion through the explosive gases, t_{arc} refers to the limitation on LI , t_{rt} is the instability growth time given by Eq.(4), and t_{mix} accounts for any additional time needed to mix the materials sufficiently. This last time is very difficult to estimate, since some mixing will occur on the "way in." At most, it should be of the order of t_{rt} . The simple estimate of t_{diff} made above

should suffice, although more detailed estimates taking account of the two-dimensional nature of the diffusion as well as the effects of the finite shunting resistance of the plasma channel and the effect of the load properties can be made.

The possibility of "seeding" the instability to enhance mixing is under investigation. The most practical way to introduce initial perturbations would be on the explosive surface. The amplitude should probably be such that the instability starts in the linear regime; this is because we want a sudden onset of mixing, and this is best achieved through the exponential linear growth of the small amplitude regime than the slower nonlinear growth. The best wavelength for Rayleigh-Taylor growth will be roughly d_{ta} or somewhat larger.¹⁵ This does not consider the role of mixing and temperature diffusion. These effects would reduce the optimal scales. Classical Spitzer heat conduction² has a diffusion coefficient $K = 5 \times 10^{21} T_{ev}^{5/2} / (Z n_e \ln \Lambda) \text{ cm}^2/\text{s}$, with (approximately) the Coulomb logarithm $\ln \Lambda = 1$, the average ionization $Z = 1$, the electron number density $n_e = 2 \times 10^{21}$, and the temperature $T_{ev} = 6 \text{ eV}$ at turn around, suggesting that for diffusion time scales of order 100 ns we need scales of order $1/10 d_{ta}$, i.e. $5 \times 10^{-3} \text{ cm}$. In the face of these uncertainties, we therefore would suggest experimental tests of seeding with initial wavelengths of order the channel thickness at turn around and initial amplitude of order of the wavelength.

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Figure Captions

1. Sketch of Pavlovskii switch.
2. Diagram of the circuit used to model the Pavloskii switch in numerical simulations.
3. Plots of gap size d (mm), current I (MA) and switch resistance R (milliohm) vs. time for the "canonical" numerical simulation of the operation of the Pavlovskii switch.
4. Selected temperature, pressure, and velocity profiles from the "canonical" numerical simulation of the switch.
 - (a) Temperature profile at $1 \mu s$. The maximum at the inner channel edge is due to Joule heating, which is concentrated at the edge because the increasing field (current) must diffuse into the channel.
 - (b) Temperature at $3 \mu s$. The channel has expanded into the air gap between the explosive and mylar, pushing the mylar inward.
 - (c) Temperature at $9.9 \mu s$. The channel has been compressed by the explosive front, being heated by the compression.
 - (d) Velocity profile at $4 \mu s$. The negative velocities are the result of inward channel expansion; the positive velocities near the wall are due to a pinch-effect generated pressure wave.
 - (e) Velocity profile at $7 \mu s$. The explosive front has reached a radius of approx. 9.3 cm. The pressure wave has reflected off the outer boundary of the channel (assumed to be a rigid wall for computational purposes).
 - (f) Velocity profile at $9.2 \mu s$. The explosive is compressing the channel.
 - (g) Velocity profile at $9.8 \mu s$. Rising internal pressure in the channel has reduced the velocity of compression. Turn around will soon occur.
 - (h) Pressure profile at $7 \mu s$. The rarefaction near 9.2 cm results from the interaction of the explosive front and the free surface of the explosive. The explosive front has reached 9.4 cm radius.
 - (i) Pressure profile at $9.8 \mu s$. The channel has been compressed and is near turn around, with pressure in the channel exceeding that in the explosive products.
5. Numerical simulation results for the channel artificially "mixed" at turn around. Times shown are after turn around.
 - (a) Temperature profile at $.2 \mu s$. The temperature is rather uniform in the channel material and has risen to 2 eV from an initial value at turn around of .6 eV.
 - (b) Pressure profile at $.2 \mu s$. The pressure is rather uniform.
 - (c) Pressure profile at $.5 \mu s$. Note the expansion of the channel, resulting in the rarefaction of the inner channel edge.
6. Results from experiments on planar analog of Pavlovskii's explosive switch, performed at Sandia.
 - (a) Resistance vs. time for experiment no. 3. The rise in switch resistance from 31-33 microseconds is much less than the peak resistance at 35 microseconds.

(b) Voltage accross switch vs. time for experiment no. 3. This figure shows the 31-33 microsecond resistance rise very clearly, as well as the later and more substantial resistance rise (as the current is falling, the change in voltage is not as dramatic as the resistance change).

(c) As in Fig. 4b, but for experiment no. 2 (no resistance plot was available for this experiment).

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TIME 2.003⁻⁷

PRESSURE (dynes/cm²)

(b) Pressure profile at .2 μ s. The pressure is rather uniform.

Figure 6. Results from experiments on planar analog of Pavlovskii's explosive switch, performed at Sandia.
(a) Resistance vs. time for experiment no. 3. The rise in switch resistance from 31-33 microseconds is much less than the peak resistance at 35 microseconds.

EXPLOSIVE SWITCH NO 3 3-13-3-80
VOLTAGE ACROSS SWITCH

(b) Voltage across switch vs. time for experiment no. 3. This figure shows the 31-33 microsecond resistance rise very clearly, as well as the later and more substantial resistance rise (as the current is falling, the change in voltage is not as dramatic as the resistance change).

VOLTAGE EXPLOSIVE SWITCH NO 2 2-12-80
ASSUMING 6 NH LOAD

(c) As in Fig. 4b, but for experiment no. 2 (no resistance plot was available for this experiment).

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