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Cutoff of the current in plasma opening switches

SERGEY V. LOGINOV

Institute of High Current Electronics SB RAS, Tomsk 634055, Russia
(loginov@oit.hcei.tsc.ru)

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Abstract. The paper considers the erosion mechanism of operation of nano- and microsecond plasma opening switches (POSs). For this purpose, postulates of the well-known erosion model of the POS operation are reviewed and some remarks on its individual statements are discussed. A voltage scaling for a nanosecond switch with rarefied plasma ($\sim 10^{13}$ cm^{-3}) is derived. It is shown that the peak voltage across the nanosecond switch on its opening is proportional to the switch conduction current. The formation of an erosion gap in a microsecond switch with high-density plasma ($\sim 10^{15}$ cm^{-3}) is put to phenomenological analysis and a voltage scaling for the switch is obtained. It is found that in the microsecond switch, unlike the nanosecond switch, the peak voltage is inversely proportional to the switch conduction current.

1. Introduction

The efficient application of plasma opening switches (POSs) for interruption of megaampere current pulses of duration $\sim$ 50 ns and energy delivery to a load with substantial pulse compression (Meger et al. 1983; Neri et al. 1987) has stimulated research to increase the energy characteristics of generators with an inductive energy store and a POS. POSs are attractive for their very simple design being a vacuum coaxial segment filled with fully ionized and, as a rule, radially injected plasma (Fig. 1). During the conduction phase, the switch separates the load from the energy store. At a certain current, the switch resistance increases sharply, thus providing fast energy delivery to the load.

The development of POS's requires solution of two main problems. The first problem is to determine the mechanism of magnetic field transport through the plasma. This mechanism determines the POS conduction time and hence the stored energy. The second problem is to ascertain the causes for current interruption and the factors responsible for the rate of energy delivery to a load. This problem is critical for prospects of the technology.

Comparative analysis of even first experiments with nanosecond POSs revealed that the increase in stored energy does not contribute to voltage multiplication and pulse compression (Weber et al. 1987). Later experiments demonstrated that with the same primary stored energy and current, a tenfold increase in the conduction time of a microsecond switch compared to that of a nanosecond switch involves a proportional increase in the time of energy delivery to the load and hence a decrease in pulse power. This negative tendency required comprehensive studies to elucidate the dynamics of the processes occurring in microsecond POSs. Numerous experiments showed that in microsecond POSs, the increasing magnetic field pressure or magnetic piston causes plasma aggregation and plasma density redistribution (Rix et al. 1991; Commisso et al. 1992), with the result that the switch current is cut off as the current channel reaches the downstream plasma edge (Grossmann et al. 1995). Because increasing the switch conduction time requires an appropriate increase in plasma density, the opening of the switch is inevitability impeded. For this reason, projection of positive experience with nanosecond POSs into the microsecond range has not lived up to expectations.

The present paper considers the erosion mechanism of operation of nano- and microsecond POSs. In Sec. 2, postulates of the well-known erosion model of the POS operation (Ottinger et al. 1984) are reviewed and some remarks on its individual statements are discussed. In Sec. 3, a voltage scaling for a nanosecond POS with rarefied plasma is derived. Section 4 examines the feasibility of current cutoff in the dense plasma of a microsecond POS and analyzes phenomenology of the formation of an erosion gap. Section 5 gives concluding remarks.

2. Erosion mechanism of the POS opening

The current cutoff in POSs is best to understand using the concept of the formation of a double space charge layer (Widner and Poukey 1976). During the formation of the layer, the extraction of ions involves motion of the anode boundary deep into the plasma. The further increase in gap width is added to by elongation of electron paths in a magnetic field. At a rather high current, electron paths are so curved that they fail to reach the anode plasma boundary. As a result, the gap resistance increases with a rate of $\sim 1\Omega$ ns^{-1}.

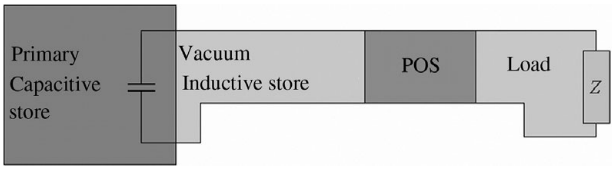

Figure 1. Schematic of pulse compression.

The double layer concept is basic for the erosion model (Ottinger et al. 1984) according to which the POS operation is divided into four successive phases: conduction, erosion, enhanced erosion, and magnetic insulation. Let us dwell briefly on these phases.

The erosion model states that during the conduction phase, the radially directed current increases uniformly over the plasma length. This postulate ignores the problem of magnetic field transport through fully ionized plasma; nevertheless, the model assumes a uniform bipolar current flow for which $I_i/I_e = (Zm/m_i)^{1/2}$, where m and m_i are the electron and ion masses, and Z is the ion charge.

The ion current is due to a radial plasma flow with a drift velocity v_d. Therefore, the conduction current in the uniform plasma is limited by the saturation current:

$$I_{\mathrm{sat}} = (m_i/Zm)^{1/2}\, 2\pi r_c l Z e n_i\, \mathrm{v}_d, \tag{1}$$

where e is the electron charge, r_c is the cathode radius, l is the plasma length, and n_i is the ion density. Saturation current (1) is proportional to the ion current density $j_{i\mathrm{sat}} = Zen_i\ \mathrm{v}_d$. It is this linear dependence of the conduction current on the plasma density, which is observed in nanosecond POSs (Weber et al. 1987).

As saturation current (1) is reached, plasma erosion begins. The ion drift can no longer provide the required current flow, and therefore, an erosion gap is formed with the rate

$$\dot{d}(t) = \{[j_i(t)/j_{i\mathrm{sat}}] - 1\}\ \mathrm{v}_d. \tag{2}$$

The current flow is bipolar as before and the voltage across the gap is determined by the Child–Langmuir law. The erosion phase ends when the electron gyroradius compares with the gap width.

In the enhanced erosion phase, as the model states, the electron paths are elongated to about the initial plasma length with attendant rapid ion extraction up to $I_i/I_e \sim (Zm/m_i)^{1/2} l/d$. The geometric factor l/d greatly increases the ion current compared to $I_i/I_e = (Zm/m_i)^{1/2}$ and assumes that the gap is formed instantaneously over the entire plasma length. This hardly probable assumption ignores the plasma dynamics during the conduction phase. It is more realistically to hypothesize that the gap develops gradually along the cathode (Fruchtman 1996; Fruchtman et al. 1999); however, this hypothesis developed in the framework of electron magnetohydrodynamics gives a conduction current scaling inconsistent with experimental data (Weingarten et al. 1999).

The enhanced ion extraction rapidly widens the gap, thus causing current cutoff and energy delivery to the load. Once the load current becomes sufficient for magnetic insulation of electrons, the final phase begins in which the only current in the gap is the space charge-limited ion current. During this phase, a further increase in gap width is possible but under magnetic field pressure due to the current switched to the load rather than under enhanced erosion.

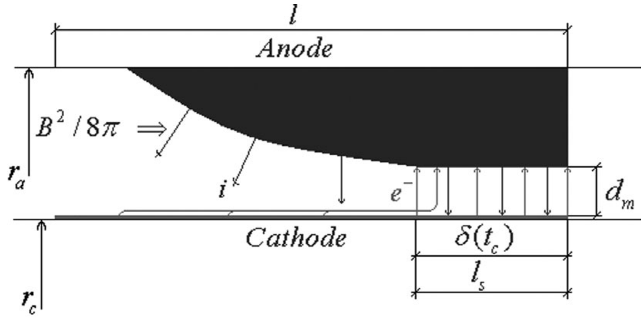

Figure 2. Formation of the gap in the regions $\delta(t_c)$ and l_s for nano- and microsecond POSs, respectively. Electrons and ion traces are marked by e^- and i.

3. Opening of nanosecond switches

A distinguishing feature of nanosecond POSs is the high rate of rise of the current ($\dot{I} \sim 10$ kA ns^{-1}) in low-density plasma ($\sim 10^{13}$ cm^{-3}). As a result, the density of the linearly increasing current $j(t) = \dot{I}t/2\pi r_c(\nu_0 t)^{1/2}$ due to magnetic field diffusion into the plasma with magnetic viscosity $\nu_0 = c^2/4\pi\sigma_0$, where c is the velocity of light, σ_0 is the initial plasma conductivity, becomes instantly higher than the saturation current density. Actually, for example for the GAMBLE II switch (Neri et al. 1987) ($\dot{I} = 15$ kA ns^{-1}, $r_c = 2.5$ cm, $l = 10$ cm), the saturation current density $j_{\mathrm{sat}} = I_c/2\pi r_c l \approx 5$ kA cm^{-2}. By diffusion in plasma with $\sigma_0 \approx 10^{14}$ s^{-1}, the current density $j(t)$ becomes equal to j_{sat} even after 20 ps. Therefore, an erosion gap in which magnetic insulation suppresses the radial electron current is formed lengthwise the plasma almost from the very beginning of the current pulse (Fig. 2; the erosion gap develops near the cathode with subsequent plasma aggregation by the magnetic field pressure in the rest part of the interelectrode gap).

For $\sigma(t) = \sigma_0 t_0/t$ (this follows from the experimental fact $j(t) \approx$ Const in the conduction phase (Weber et al. 1987)), integration of (2) with the knowingly valid condition $\dot{d}(t) \gg \mathrm{v}_d$ gives the time $t_0 = (2\pi^{1/2}(\mathrm{v}_{T_e}/\mathrm{v}_d)(mc/e)\{j_{\mathrm{sat}}/[\sigma_0^{1/2}(2\dot{I}/cr_c)^2]\})^{2/5}$ during which the width of the erosion gap compares with the Larmor radius $r_L(t_0) = \mathrm{v}_{T_e}/[e(2\dot{I}t_0/cr_c)/mc]$, where v_{T_e} is the electron thermal velocity. Because of the erosion, the conduction time is $t_c \sim [l/2r_L(t_0)]t_0$ (for $\mathrm{v}_{T_e} = 0.2$ cm ns^{-1}, $\mathrm{v}_d = 10$ cm μs^{-1} the estimation gives $t_0 \approx 0.4$ ns, $r_L(t_0) \approx 0.2$ mm, $t_c \sim 100$ ns), the switch conduction current is determined by saturation current (1).

On termination of the conduction phase, the current density in a channel of width $\delta(t_c) \sim [\nu_m(t_c)t_c]^{1/2}$ is

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higher than the saturation current density by a factor of $\sim l/\delta(t_c)$. Then, from relation (2) follows $d(\tau) = l\mathrm{v}_d\tau/[v_m(t_c)t_c]^{1/2}$. Hereinafter, the time $\tau = t - t_c$ is counted from the instant at which the conduction phase terminates.

The electron current density is given by the Child–Langmuir law, which in the ultrarelativistic case is expressed as

$$j_e(\tau) = (mc^3/e)[2\pi d^2(\tau)]^{-1}[eU_s(\tau)/mc^2], \tag{3}$$

where $U_s(\tau)$ is the voltage across the forming gap.

Accordingly to the erosion model (Ottinger et al. 1984), the switch current is cut off when the gap at the point in time τ_s compares with the electron gyroradius equal (as it follows from formula (13a) at $\gamma \gg 1$) to

$$d_m(\tau_s) = [(mc^3/2e)/I_c][eU_s(\tau_s)/mc^2]r_c, \tag{4}$$

where the switch conduction current is $I_c \approx 2\pi r_c[v_m(t_c)t_c]^{1/2}j_e(\tau_s)$, assuming that during the time τ_s the decrease in current is negligible and $j_e(\tau_s) \approx j_e(t_c)$.

From (3) and (4), it follows that $d_m(\tau_s) = 2[v_m(t_c)t_c]^{1/2}$, where $\tau_s = 2v_m(t_c)t_c/l\mathrm{v}_d$ is the gap formation time. For $t_c \approx 50$ ns, the current channel width $[v_m(t_c)t_c]^{1/2} \sim 2$ cm, and then $\tau_s \approx 80$ ns, which is unreally high for nanosecond switches. This means that the decay of the switch current on opening should be taken into account. For this purpose, it is quite possible to represent the switch current $I_s(\tau)$ and the generator discharge current $I_g(\tau)$ in the form $I_s(\tau) = I_c\exp(-\tau/\tau_{cs})$ and $I_g(\tau) = I_c\exp(-\tau/\tau_{cg})$, where τ_{cs} and τ_{cg} are the characteristic times of the electrical circuit ($\tau_{cg} \geqslant \tau_{cs}$). In this case, integration of (2) gives $d(\tau) = l\mathrm{v}_d\tau_{cs}[1 - \exp(-\tau/\tau_{cs})]/[v_m(t_c)t_c]^{0.5}$.

As shown above, the electrons become magnetized very quickly (the time t_0 is very small); therefore, the discharge current $I_g(\tau)$ coincides with magnetic insulation current (13a). Consequently, the switch voltage can be calculated from the equation $I_g(\tau) = (mc^3/2e)[eU_s(\tau)/mc^2]r_c/d(\tau)$ for $\gamma \gg 1$. Assuming for simplicity that $\tau_{cg} = \tau_{cs}$ (for differing τ_{cs} and τ_{cg}, the analysis is easy to summarize), this equation allows us to determine the switch peak voltage

$$eU_s(\tau_m)/mc^2 = [I_c/(2mc^3/e)]l\mathrm{v}_d\tau_{cs}/\{[(v_m(t_c)t_c]^{1/2}r_c\}, \tag{5}$$

where $\tau_m = \tau_{cs}\ln 2$.

The voltage (5) is proportional to the conduction current. This agrees with the experiments: as I_c was increased from ∼250 kA (Meger et al. 1983) to ∼750 kA (Neri et al. 1987), the peak voltage increased from ∼1.4 MV to ∼4.3 MV (in these experiments, $[v_m(t_c)t_c]^{1/2}$ and r_c were found approximately equal). Moreover, formula (5) yields the reasonable numerical result: $U_s(\tau_m) = 2.2 - 4.4$ MV for $\tau_{cs} = 10 - 20$ ns. For further quantitative comparison with experimental current traces, we can use the rate of rise of the switch resistance $\dot{R}_s = [U_s(\tau_m)/I_s(\tau_m)]/\tau_m = (\mathrm{v}_d/cr_c 2\ln 2)l/[(v_m(t_c)t_c]^{1/2}$, which does not depend on τ_{cs} (for the discussed experiments, $\dot{R}_s \approx 0.4\Omega$ ns^{-1}).

4. Voltage scaling for microsecond switches

In microsecond switches, unlike nanosecond switches, the conduction current is an order of magnitude lower than bipolar current (1) and varies according to the scaling (Rix et al. 1991; Commisso et al. 1992):

$$I_c = (12\pi n_i m_i c^2)^{1/4}(\dot{I}r_c l)^{1/2}, \tag{6}$$

which points to plasma aggregation by a magnetic piston. Comparison of (6) with (1) suggests that the condition $I_c \ll I_{\mathrm{sat}}$ is knowingly valid for high-density plasma with $n_i \sim 10^{15}$ cm^{-3}.

The different current scaling, in this case, owes to a different mechanism of magnetic field penetration. Because in a microsecond switch, the plasma density is two orders of magnitude higher and the rate of rise of the current is an order of magnitude lower compared to a nanosecond switch, the microsecond switch reveals no prerequisite for plasma erosion up to the point in time at which the current channel reaches the downstream plasma edge (Loginov 2011b). On termination of the conduction phase, the axial plasma length accessible for current passage is reduced to l_s for which the saturation current is

$$I_{\mathrm{sat}l_s} = (m_i/Zm)^{1/2}2\pi r_c l_s Ze\tilde{n}_i\,\mathrm{v}_d. \tag{7}$$

Saturation current (7) depends linearly on the density $\tilde{n}_i$ at the downstream plasma edge. For perfect 1D plasma aggregation, this current coincides with (1) because the plasma mass with no loss remains unchanged. This means that $\tilde{n}_i l_s \approx n_i l$ without regard for details of the density redistribution. The condition $I_c \ll I_{\mathrm{sat}}$ is thus valid as before and no prerequisite for plasma erosion arises. Hence, there should be further acceleration of the plasma compressed in the current channel. Actually, this does not take place for two reasons. The first reason is that even 1D compression involves leakage of the piston due to finite conductivity of the current channel. The second and more significant reason is that the motion of the plasma in the conduction phase is 2D, which is particularly pronounced in coaxial POSs with large interelectrode gaps. The radius-dependent magnetic field pressure causes radial pushing of the plasma toward the electrodes (the plasma boundary facing the energy store becomes similar to that shown in Fig. 2). Because of the above processes, the current $I_{\mathrm{sat}l_s}$ can be much less than I_{sat}. Another fact assisting this decrease is that in actual switches, the current channel on termination of the conduction phase enters the region in which the density decreases toward the load. If for all the above reasons, $\tilde{n}_i$ approximates n_i, the switch length l_s required for current passage decreases to a value much less than l. The assumption of $\tilde{n}_i \sim n_i$ is confirmed by the numerical simulation: the increase in density after aggregation of even homogeneous plasma is no greater than about twofold (Xu and Wang 2006).

Let us consider phenomenology of the formation of an erosion gap in plasma of length l_s with a certain average density $\tilde{n}_i$ neglecting its possible variation during the short opening time. The length l_s is defined as $l_s = (I_c/I_{sat})l/\xi$, where the saturation current I_{sat} is given by (1), switch current (6) coincides with (7), and $\xi = \tilde{n}_i/n_i$. Increasing the plasma density decreases $l_s \propto \xi^{-1} n_i^{-3/4}$.

Further, it is possible to assume that l_s decreases with a certain rate $\tilde{u}(t)$. The opening time is short compared to the conduction time and, hence, $l_s(t) = l_s - \int_{t_c}^{t} \tilde{u}(t')dt' \approx l_s - \tilde{u}(t_c)(t - t_c)$; from whence it follows that the opening time is $\tau_s = t - t_c = l_s/\tilde{u}(t_c)$. With small values of τ_s, we can also neglect the downstream displacement of the current channel on opening. If l_s decreases with a rate $\tilde{u}(t_c) = \delta_s u_{mp}(t_c)$, where the coefficient δ_s takes into account the difference between the magnetic piston velocity $u_{mp}(t_c)$ and the rate of decrease of l_s due to axial plasma motion, the time τ_s in view of scaling (6) is equal to

$$\tau_s = 3(mc^2/e)(m_i/Zm)^{1/2} r_c l\dot{I}/\xi\delta_s \, \mathrm{v}_d I_c^2. \tag{8}$$

For any coordinate z in the range $0 \leqslant z \leqslant l_s$ during the time $\tau_1 = z/\tilde{u}(t_c)$, the expansion of the gap due to the decrease in l_s obeys the equation:

$$\dot{d}(\tau) = \left(\{l_s/[l_s - \tilde{u}(t_c)\tau]\} - 1\right) \mathrm{v}_d. \tag{9a}$$

Once the electron current is cut off, the expansion rate in a time $\tau_2 = (l_s - z)/\tilde{u}(t_c)$ is determined by the excess of the extracted ion current over the saturation current:

$$\dot{d}(\tau) = \{[j_i(\tau)/j_{isat}] - 1\} \, \mathrm{v}_d. \tag{9b}$$

Here, $j_i(\tau) = \mu(Zm/m_i)^{1/2}(\sqrt{2}/9\pi)(mc^3/e)[eU_s(\tau)/mc^2]^{3/2} d^{-2}(\tau)$ is the ion current density, μ is the factor of possible enhancement of the extracted ion current against the background of the space charge of magnetized electrons (Bergeron 1976).

Integration of (9) together with circuit equations gives the gap profile along l_s. For analytical solution, one can assume that on opening the switch current decreases slightly. This assumption is not quite correct, because the current decay depends on the load parameters; however, this affects only the numerical factor in the voltage scaling, leaving its functional tendencies unchanged.

From (9a) follows the dependence

$$\tilde{d}(z) = [\mathrm{v}_d/\tilde{u}(t_c)]\{l_s \ln[l_s/(l_s - z)] - z\},$$

which reflects the logarithmically sharp increase in $\tilde{d}(z)$ at $z \to l_s$.

Assuming that the voltage across the gap rises linearly $U_s(\tau) = \dot{U}_s\tau$, equation (9b) for the case $\dot{d}(\tau) \gg \mathrm{v}_d$ of practical interest is reduced to

$$d^2(\tau)\dot{d}(\tau) = b\tau^{3/2},$$

where $b = (2\sqrt{2}/9)\mu[(mc^3/e)/I_c](e\dot{U}_s/mc^2)^{3/2} \, \mathrm{v}_d r_c l_s$. Integration of this equation gives the gap profile:

$$d(z) = (\tilde{d}^3(z) + (6b/5)\{[l_s/\tilde{u}(t_c)]^{5/2} - [z/\tilde{u}(t_c)]^{5/2}\})^{1/3}. \tag{10}$$

Unlike $\tilde{d}(z)$, the second term in (10) decreases with an increase in z and the gap profile is thus non-monotonic. The minimum value of $d(z)$ is reached at the coordinate $z_m = l_s - \Delta$, which is determined from the equation:

$$\Delta \ln^{-2}(l_s/\Delta \, \mathrm{e}) = \mathrm{v}_d^3 l_s^{3/2}/b\tilde{u}^{1/2}(t_c). \tag{11}$$

Substitution of (11) in (10) gives the minimum gap

$$d_m = \tilde{\Lambda} \, \mathrm{v}_d \tau_s, \tag{12}$$

where $\tilde{\Lambda} = [1 + 3/\ln(l_s/\Delta \, \mathrm{e})]^{1/3} \ln(l_s/\Delta \, \mathrm{e})$. Hereinafter, the symbol e is the base of a natural logarithm.

The current is cut off when the formed gap becomes magnetically insulated. In the one-particle approximation, this condition is fulfilled for the current

$$I_c \approx I_{mi} = (mc^3/2e)(\gamma^2 - 1)^{1/2} r_c/d_m, \tag{13a}$$

where $\gamma = 1 + eU_s(\tau_s)/mc^2$. At voltages $eU_s(\tau_s)/mc^2 \ll 1$, formula (13a) is reduced to

$$I_{mi} \approx (mc^3/\sqrt{2}e)[eU_s(\tau_s)/mc^2]^{1/2} r_c/d_m. \tag{13b}$$

Criterion (13) is automatically valid throughout the time of intense ion extraction (Black et al. 2000). Indeed, increasing d_m decreases the current required for magnetic insulation. If I_{mi} becomes less than the switch current, the electron sheath is pressed to the cathode plasma. Thus, the cause for ion current enhancement disappears. The decrease in extracted ion current to the space charge-limited current slows down the gap expansion. If, on the contrary, the gap decreases for some reason, the required current I_{mi} becomes higher than the switch current. Straightening the electron paths increases the extracted ion current and the expansion of the gap is resumed. Thus, the gap cannot differ greatly from that required for magnetic insulation.

Substitution of (12), (13) in (11) excludes the dependence on the rate of rise of the voltage $\dot{U}_s = U_s(\tau_s)/\tau_s$ and gives transcendental equations for the coordinate at which the gap has a minimum width. At high voltages ($\gamma \gg 1$), the coordinate Δ is determined from the equation:

$$(l_s/\Delta)^2 \; [3 + \ln(l_s/\Delta \, \mathrm{e})]^{-1} \; \ln^2(l_s/\Delta \, \mathrm{e}) = (64\mu^2/81)[I_c/(mc^3/e)](l_s/r_c)[\tilde{u}(t_c)/\mathrm{v}_d].$$

At low voltages ($\gamma \sim 1$), Δ is determined by the equation:

$$\Delta \ln[1 + 3/\ln(l_s/\Delta \, \mathrm{e})] = (9/8\mu)[(mc^3/e)/I_c]^2 (r_c^2/l_s)[\tilde{u}(t_c)/\mathrm{v}_d].$$

With (8) and current scaling (6), expression (12) is reduced to the very simple form:

$$d_m = (\sqrt{3}\tilde{\Lambda}/\xi\delta_s)c/\omega_{pe}, \tag{14}$$

where $\omega_{pe} = (4\pi ne^2/m)^{1/2}$ is the electron plasma frequency and $n = Zn_i$. The increase in plasma density decreases $d_m \propto n^{-1/2}$. This tendency is confirmed (accurate to a power exponent) by the numerical simulation of the gap formation in plasma with $n = 10^{14}$–5×10^{15} cm^{-3} (Grossmann et al. 1995). The additional increase in

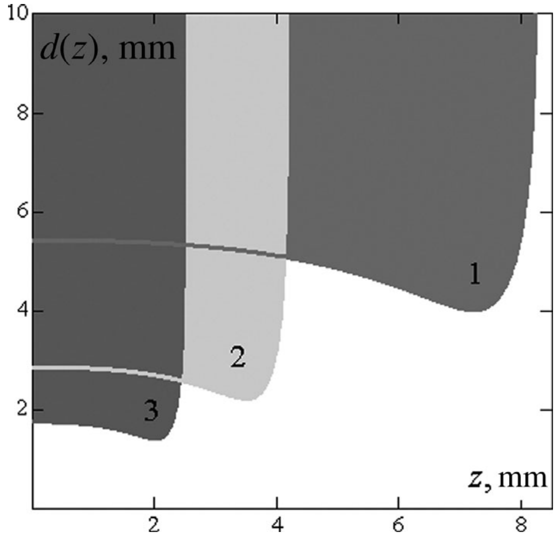

Figure 3. Gap profile along l_s at $n(10^{15}\ \mathrm{cm}^{-3}) = 1$ (*1*), 2.5 (*2*), and 5 (*3*).

plasma density at the downstream plasma edge after aggregation impedes the plasma erosion because $d_m \propto \xi^{-1}$.

Figure 3 shows the gap profile in relation to the density (l =13 cm, r_c =3.8 cm, $\dot{I}$ =1.5 kA ns^{-1}, v_d =2 cm μs^{-1}; for doubly ionized carbon plasma, $Z = 2$, $(m_i/Zm)^{1/2} = 105$; $\mu = 1$). The increase in density, first, decreases l_s and, second, decreases d_m. The dependence of d_m on the factor μ is logarithmically weak and variations of μ affect only the axial position of d_m. In all cases, it is put for simplicity that $\xi = 1$ (in real switches, $\tilde{n}_i$ can presumably be taken as an average of n_i over the plasma length) and $\delta_s \approx 0.13$ (the plasma velocity at the front of the current channel is $\sim 0.87u_{mp}(t_c)$ (Loginov 2011b)). A negative influence of the increase in ξ and δ_s on d_m is obvious from formula (14).

From (13), in view of (8) and (12), follows the voltage scaling for the magnetically insulated gap. At $\gamma \gg 1$, the scaling is written as

$$eU_s(\tau_s)/mc^2 \approx 6(m_i/Zm)^{1/2}(\tilde{\Lambda}/\xi\delta_s)(l\dot{I}/cI_c). \tag{15a}$$

At $\gamma \sim 1$, the scaling is given by the different formula:

$$eU_s(\tau_s)/mc^2 \approx 18(m_i/Zm)(\tilde{\Lambda}/\xi\delta_s)^2(l\dot{I}/cI_c)^2. \tag{15b}$$

Directly from (13a), it follows that at $\gamma \gg 1$, the switch resistance $R_s[\Omega] = 60d_m/r_c$ is proportional to the minimum gap width. For the gap formation time $\tau_s = l_s/\tilde{u}(t_c)$, the rate of rise of the resistance $\dot{R}_s[\Omega\ /\mathrm{s}] = 60\tilde{\Lambda}\ \mathrm{v}_d/r_c$ depends mainly on the cathode radius and plasma drift velocity. The first result is clear because maintaining the extracted ion current with a decrease in r_c requires faster expansion of the gap. The second one is of no practical significance because the attainable resistance does not depend on the plasma drift velocity to the cathode. With the gap profiles shown in Fig. 3 (they were calculated for the open-circuit regime of POS operation), the rate of rise of the resistance $\dot{R}_s$ is no greater than $\sim$ 100 mΩ ns^{-1} (Fig. 4). Increasing the density decreases $\dot{R}_s$, whereas enhancing the ion extraction at varying μ increases $\dot{R}_s$. If the cathode radius is increased and a low-inductive load is turned on, $\dot{R}_s$ decreases down to less than tens of milliohm per nanosecond (Loginov 2011a).

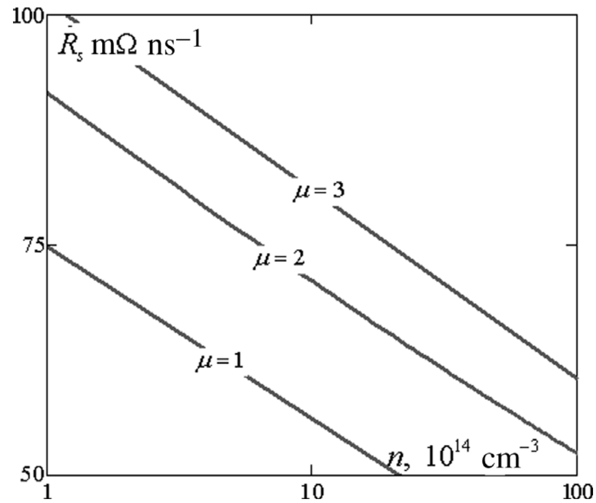

Figure 4. Dependence of $\dot{R}_s$ on n at varying μ.

From $\dot{R}_s[\Omega\ /\mathrm{s}] = 60\tilde{\Lambda}\ \mathrm{v}_d/r_c$, it follows that the dependence of $\dot{R}_s$ on the switch current and degree of plasma compression is logarithmically weak. However, this conclusion is valid only at $\gamma \gg 1$. At $\gamma \sim 1$, the resistance is given by another formula: $R_s[\Omega] = 60[I_c/(mc^3/e)](d_m/r_c)^2$. Because $d_m \propto \xi^{-1}$, the increase in density at the downstream plasma edge decreases $R_s \propto \xi^{-2}$ and $\dot{R}_s \propto \xi^{-1}$. Additionally, $\dot{R}_s$ becomes dependent on the conduction current: $\dot{R}_s \propto l\dot{I}/I_c$.

For high voltages, which are of practical interest, the scaling $eU_s(\tau_s)/mc^2 \propto l\dot{I}/I_c$ is confirmed by the experimentally observed regularities of current cutoff in the microsecond megaampere switches (Loginov 2009, 2011a). At the same time, the inversely proportional dependence of the switch voltage on the conduction current means that the power dissipated in the switch is proportional to the product of the initial plasma length into the rate of rise of the current. Because the conduction current is limited by the short-circuit current I_0, the maximum length cannot be greater than $l \leqslant (12\pi m_i n_i c^2)^{-1/2} I_0^2/r_c\dot{I}$. Hence, the peak power varies according to $P_m \propto I_0^2/r_c n^{1/2}$ or $P_m \propto U_0^2/r_c n^{1/2}$ (here, $I_0 = U_0/\rho$ with U_0 being the output voltage of the primary capacitive store and ρ being the characteristic impedance of the discharge circuit). With a fixed plasma density, the power can be increased either by decreasing the cathode radius or by increasing the output voltage of the primary store. Increasing the plasma density decreases the peak power.

5. Conclusion

The inversely proportional dependence of the switch peak voltage on the conduction current radically distinguishes microsecond POSs from nanosecond ones. For the latter, the peak voltage is given by scaling (5), which depends linearly on the conduction current.

Scaling (15a) means that the problems of increasing the conduction time and attaining high resistance on current cutoff are not only interrelated but they are also contradictory. This is a natural outcome of the necessity to increase the plasma density to provide the desired current with a microsecond rise time. In nanosecond switches, the erosion of the rarefied plasma begins almost from the very beginning of the rapidly rising current pulse, whereas in microsecond switches the tenfold decrease in the rate of rise of the current and the increase in plasma density precludes the erosion throughout the conduction phase and reduces the high-density plasma dynamics to aggregation. As a result, the conditions for the onset of erosion are attainable only when the current channel reaches the downstream edge of the initial plasma and the available plasma length decreases to a value insufficient for bipolar current flow.

References

Bergeron, K. D. 1976 Two-species flow in relativistic diodes near the critical field for magnetic insulation. *Appl. Phys. Lett.* **28**(6), 306–308.

Black, D. C., Commisso, R. J., Ottinger, P. F., Swanekamp, S. B. and Weber, B. V. 2000 Experimental determination of gap scaling in a plasma opening switch. *Phys. Plasmas* **7**(9), 3790–3796.

Commisso, R. J., Goodrich, P. J., Grossmann, J. M., Hinshelwood, D. D., Ottinger, P. F. and Weber, B. V. 1992 Characterization of a microsecond-conduction time plasma opening switch. *Phys. Fluids* B **4**(7), 2368–2376.

Fruchtman, A. 1996 A vacuum sheath propagation along a cathode. *Phys. Plasmas* **3**(8), 3111–3115.

Fruchtman, A., Grossmann, J. M., Swanekamp, S. B. and Ottinger, P. F. 1999 Sheath propagation along the cathode of a plasma opening switch. *IEEE Trans. Plasma. Sci.* **27**(5), 1464–1467.

Grossmann, J. M., Swanecamp, S. B., Ottinger, P. F., Commisso, R. J., Hinshelwood, D. D. and Weber, B. V. 1995 Gap formation processes in a high-density plasma opening switch. *Phys. Plasmas* **2**(1), 299–309.

Loginov, S. V. 2009 Plasma opening switch operation scenario. *IEEE Trans. Plasma Sci.* **37**(10), 1930–1935.

Loginov, S. V. 2011a Electron loss in microsecond megaampere plasma opening switches. *IEEE Trans. Plasma Sci.* **39**(12), 3386–3390.

Loginov, S. V. 2011b Plasma dynamics in microsecond megaampere plasma opening switches. *Phys. Plasmas* **18**(10), 102104 (6 pp.).

Meger, R. A., Commisso, R. J., Cooperstein, G. and Goldstein, S. A. 1983 Vacuum inductive store/pulse compression experiments on a high power accelerator using plasma opening switches. *Appl. Phys. Lett.* **42**(11), 943–945.

Neri, J. M., Boller, J. R., Ottinger, P. F., Weber, B. V. and Young, F. C. 1987 High-voltage, high-power operation of the plasma erosion opening switch. *Appl. Phys. Lett.* **50**(19), 1331–1333.

Ottinger, P. F., Goldstein, S. A. and Meger, R. A. 1984 Theoretical modeling of the plasma erosion opening switch for inductive storage applications. *J. Appl. Phys.* **56**(3), 774–784.

Rix, W., Parks, D., Shanon, J., Thompson, J. and Waisman, E. 1991 Operation and empirical modeling of the plasma opening switch. *IEEE Trans. Plasma Sci.* **19**(2), 400–407.

Weber, B. V., Commisso, R. J., Cooperstein, G., Grossmann, J. M., Hinshelwood, D. D., Mosher, D., Neri, J. M., Ottinger, P. F. and Stephanakis, S. J. 1987 Plasma erosion opening switch research at NRL. *IEEE Trans. Plasma Sci.* **15**(6), 635–648.

Weingarten, A., Bernshtam, V. A., Fruchtman, A., Grabowski, C., Krasik, Y. E. and Maron, Y. 1999 Study of the effects of the prefilled-plasma parameters on the operation of a short-conduction plasma opening switch. *IEEE Trans. Plasma Sci.* **27**(6), 1596–1605.

Widner, M. M. and Poukey, J. W. 1976 Ion sheath motion in plasma-filled diodes. *Phys. Fluids* **19**(11), 1838–1840.

Xu, X. and Wang, Y.-N. 2006 Magnetohydrodynamic simulation of a coaxial high-density plasma opening switch. *IEEE Trans. Plasma Sci.* **34**(4), 1529–1535.