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Final Report

ACCELERATOR RESEARCH STUDIES

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SUMMARY

The program consisted of the following three tasks: TASK A, "Study of Transport and Longitudinal Compression of Intense, High-Brightness Beams," (P.I., M. Reiser); TASK B, "Study of Collective Ion Acceleration by Intense Electron Beams and Pseudospark Produced High Brightness Electron Beams", (Co-P.I.'s, W.W. Destler, M. Reiser, M.J. Rhee, and C.D. Striffler); TASK C, "Study of a Gyroklystron High-Power Microwave Source for Linear Colliders" (Co-P.I.'s, V.L. Granatstein, W. Lawson, M. Reiser, and C.D. Striffler).

The major highlights of TASK A were: 1) completion of the Hughes electron gun (5 kV, 240 mA) emittance measurements including the evaluation of errors due to space charge; 2) significant progress in obtaining agreement between simulation and measurement for the multiple-beam experiment; 3) development of a new theoretical model for predicting emittance growth in mismatched and off-centered high-brightness beams; 4) modification of the new gridded electron gun for the pulse compression experiment, measurements of the electron beam emittance, and testing of the induction gap.

Highlights in TASK B were: 1) construction of a 100 cm long laser-controlled collective ion acceleration experiment and completion of preliminary tests for two acceleration gradients; 2) further progress in our theoretical electron beam work supporting the collective ion acceleration experiments; 3) successful studies of high-brightness electron beam production at 50 kV (versus 25 kV before) in the pseudospark experiment.

Highlights in TASK C were: 1) the achievement of 3 MW of beam power in the two-cavity gyrokystron configuration which represents an increase of a factor of 50 in the state of the art; 2) completion of design and construction of the components for the three-cavity amplification experiment scheduled to get under way in January 1991; 3) significant progress in code development in support of the experiment; 4) experimental and theoretical investigation of phase variation due to voltage changes in gyrokystrons showing that this effect is dramatically different from the variation in klystrons; 5) a collider scenario study (SAIC collaboration) showing that a 100 MW gyrokystron power source at 20 GHz would provide the most economic solution for a collider in the 1-3 TeV (CM) range.

I. Introduction and Synopsis

TASK A of this contract was concerned with studies of the physics of high-brightness particle beams for advanced accelerator applications. Specifically, we were investigating several major effects that caused emittance growth and brightness deterioration such as nonlinearities in either the applied or the space charge forces and instabilities. An important feature of this work was the fact that the experiments were compared with theory and with particle simulation studies by Irving Haber and Helen Rudd at NRL/Berkeley Scholars. This comparison provided an excellent check for the validity and accuracy of the simulation codes and thereby established code credibility for design studies in areas where experimental testing is either too costly or not available for other reasons.

Prior to 1988, our research concentrated on transverse beam physics, in particular on understanding the limits and constraints affecting beam quality in long periodic focusing channels. Most of this work was carried out with our 5-keV, 240-mA electron gun developed originally with the collaboration of W. Herrmannsfeldt at SLAC and built by the Hughes Company in California. The beam from this gun is injected into a 5-meter long focusing channel built in-house and consisting of 36 periodically spaced solenoid lenses and two matching lenses.

In 1988/89 we designed and built a new injector with a gridded electron gun and an induction linac module. This injector produced short pulses with a velocity shear that allowed us to study the beam physics of longitudinal compression. During the past year the construction of this new injector and testing of the various components (electron gun, three matching lenses, induction linac), all designed and built in-house, was completed. The induction linac was designed to produce an energy shear (2.5 keV at beam front to 7.5 keV at rear end) that resulted in longitudinal compression of the 50 ns initial pulse by a factor of 3 to 5 at the end of the 5-m long focusing channel.

In Part II below, the progress made during the past 12-month period is summarized. Section A describes the results of measurements of the Hughes gun beam emittance in a special test stand. The effects of space charge and slit geometry on the emittance measurement were analyzed and, for the first time, we have a very detailed description of the beam produced by this gun. These measurements were triggered by discrepancies between simulation and experimental data in our multiple-beam experiment, as described in Section B. Our original simulation studies for the multiple beam experiment, as reported in the 1988 Phys. Rev. Letters article, showed good agreement on the beamlet merging process, but not on the image formation. New simulation studies using an initial emittance two-times larger than was first assumed, and a 10% higher magnetic field than that predicted from our

solenoidal lens data produced remarkably good agreement with the phosphor screen pictures. Our efforts during the past year to explain these higher numbers were partly successful, but some differences still remain. Thus, according to the Hughes gun measurements, the initial emittance should be somewhat higher than the factor two used in the simulation. At the same time, we have not yet been able to fully account for the 10% higher magnetic field.

Progress on the various components of the pulse compression experiment is described in Section C. A major time-consuming job that caused some delays in our schedule was the replacement of the cathode and small modifications to the electron gun. The old oxide cathode was replaced by a dispenser cathode (we bought five of these new cathodes from EIMAC/VARIAN). Dispenser cathodes can be used again after breaking the vacuum for making required changes while oxide cathodes must be discarded. We have repeated pepperpot measurements to obtain emittance data for the modified gun with dispenser cathode.

One of the most important scientific achievements of our research program on high-brightness beams has been the fact that it stimulated the development of a new theory capable of explaining effects that were previously not understood. The major example is the theory of emittance growth in beams with nonuniform charge density (Reiser, Struckmeier, and Wangler, 1983/84). This theory has since been confirmed in numerous particle simulation studies and in our recent multiple-beam experiment. During the past year the framework of this theory has been significantly broadened and generalized to include effects other than space charge nonuniformity, such as beam mismatch and off-centering. This new theory is discussed briefly in Section D, and a paper entitled "Free Energy and Emittance Growth in Nonstationary Charged Particle Beams" has been submitted for publication to *Physical Review Letters*.

II. Research Progress

A. Measurements of Hughes Gun Characteristics

Measurement of initial beam conditions such as beam expansion, current density profiles, and emittance are essential for deriving accurate theoretical and numerical models. A small diagnostic chamber for measuring these initial gun characteristics was designed and built last year. The emittance meter used is of the slit/pinhole type. Due to space charge effects in the measurement, the initial emittance data yielded inaccurate results. The value that was determined was too high (116 mm-mrad). The quantitative effect of the space charge can be estimated and is presented later. For those first measurements, the compensated emittance was calculated to be approximately 98 mm-mrad (intrinsic emittance of 85 mm-mrad, and corresponding growth of 1.15). In order to test this method of space charge compensation and to get a trustworthy experimental value of emittance, a series of slits of various widths were made and placed in the test stand chamber. By performing measurements with each of these slits and linearly fitting the emittance values as a function of slit width, the zero slit width (and hence zero space charge) emittance can be extrapolated. Finite geometry effects also have an effect on measured emittance. This method of multiple slits has been used by others¹ to compensate for finite slit effects. We will later show that the geometrical effects in this experiment are small enough to be ignored.

A diagram of the emittance measurement chamber is shown in Fig. 1. A pinhole/charge collector located 45 mm from the anode measures the current density profile. A plot of the normalized current density of a 5 kV, 245 mA full beam 45 mm from the anode is shown in Fig. 2. Also included in the figure is the profile generated by the Herrmannsfeldt electron gun simulation code, EGN.

The four vertical slits mentioned above are located 45 mm from the anode. A horizontally sweeping charge collector (pinhole/Faraday cup) 135 mm from the anode allows the emittance to be measured via the slit-pinhole method. The slits are constructed from 0.05 mm thick tantalum foil with slit widths $2X_0 = 0.254$ mm, 0.204 mm, 0.138 mm, and 0.055 mm. In the collector plane, which is a distance $L = 85$ mm from the slit, a Faraday cup assembly with a pinhole of diameter $2r_p = 0.1$ mm scans across the beam in the x -direction. The cup is mounted on an XYZ manipulator which, in turn, is driven by a computer-controlled stepper motor. Figure 3 shows in schematic form a beam profile measured with the Faraday cup pinhole when the slit is located at a distance x_s from the beam axis and the distribution is Gaussian in transverse angles $x' = \xi/L$. The mean angle $x' = \eta$ is defined by the peak of the curve while the rms width of the distribution is measured by the parameter α as shown in the figure.

Figure 1: Schematic of the test stand for measuring initial gun emittance.

Figure 2: Normalized current density profile of full 245 mA, 5 kV beam taken 45 mm from anode and corresponding Herrmannsfeldt code profile.

Figure 3: Schematic of the slit-pinhole method for emittance measurement.

An Apple IIe controls the stepper motor which employs independent closed-loop feedback to accurately drive the pinhole/cup assembly to within 0.02 mm accuracy. The signal from the collector is then displayed on a Tektronix 2430 digital scope. The Apple, via the GPIB bus, then reads and stores the current density for the present pinhole position before proceeding to the next pinhole position. The functions $\alpha(r)$, $\eta(r)$, and $n(r)$ defined below can then be determined from the individual beam profiles. The normalized current density is represented by $n(r)$.

Following Rhee and Schneider², the beam distribution function $f_4(x, x', y, y')$ can be written as

$$f_4 = \frac{n(r)}{\sqrt{2\pi}\alpha(r)} \exp\left(\frac{-(x' - \eta(r)x/r)^2 - (y' - \eta(r)y/r)^2}{2\alpha^2(r)}\right). \quad (1)$$

We obtain $f_2(x, x')$ by numerically integrating $f_4(x, x', y, y')$ over y and y' . We can then calculate $\langle x^2 \rangle$, $\langle x'^2 \rangle$, and $\langle xx' \rangle^2$ and the effective emittance of the beam which is defined as

$$\epsilon = 4\epsilon_{rms} = 4 \left(\langle x^2 \rangle \langle x'^2 \rangle - \langle xx' \rangle^2 \right)^{1/2}. \quad (2)$$

This analysis assumes that space charge and finite geometry effects are negligible in the expansion of the sheet beamlets. If these sources of error are insignificant, each slit measurement should result in the same emittance. If they are significant, the zero slit width emittance can be extrapolated from the data. In Fig. 4, the measured emittances are shown plotted as a function of slit width along with a best linear fit (upper curve in the figure) to

Figure 4: Plot of emittance vs slit width for the measured values, the space charge compensated values using the 1-dimensional sheet beam, and K-V beam models.

the points. The other curves will be explained later. The nonuniformity of the measured values is obvious. Thus, at least one of the above errors is present. The linearity of the data suggests that only one of the error types dominates. Analysis can approximately compensate for these effects. Before doing that, however, note that the most reliable value of the emittance is the extrapolated zero slit width emittance. In this case, it is 97 mm-mrad. This corresponds to a growth over the intrinsic emittance (86 mm-mrad) of 1.13.

1. Finite Slit Effects

The slit/pinhole technique and the subsequent data analysis described earlier assumes that the expansion of each beamlet between the slit and pinhole depends only on the velocity distribution of the particles. In a well designed emittance meter this is a good approximation. However, unavoidable errors are introduced by the finite size of the slit and pinhole. In addition, a finite size sheet beam also has a finite current. This space charge causes additional expansion over that due to emittance. Thus, the measured distribution at the pinhole depends not only on the velocity distribution, but on the slit width, pinhole size, and the current density of the sheet beam. These three sources of error cause the measured emittance to be higher than the actual emittance. Careful analysis must be performed to compensate for these effects.

The sheet beam can be modeled in several ways. First, a 1-dimensional model can be used when current and emittance are nonzero. The differential equation describing the envelope of the sheet beam X_m (here, X_m is the total half-width of the sheet, as opposed to the rms half width), is given by

$$X_m'' - \frac{\pi K}{2R_s} - \frac{\epsilon_x^2}{X_m^3} = 0, \quad (3)$$

where R_s is the radius of the full beam at the slit, K is the generalized perveance given by

$$K = \frac{2I_s}{I_0(\beta\gamma)^3}, \quad (4)$$

where I_s is the current of the sheet beam, $I_0 = 17,000$ A, β is v/c , and γ is the relativistic energy factor. ϵ_x is the total emittance of the sheet beam described by

$$\epsilon_x = X_0 \sqrt{\frac{3kT_s}{2qV_0}}, \quad (5)$$

where k is Boltzmann's constant, q is the charge of the particle, V_0 is the accelerating voltage, X_0 is the half slit width, and T_s is the temperature of the beam at the slit. This correctly characterizes the initial beam as a sheet, but assumes an infinite long axis. This equation has the distinct advantage of having analytic solutions when either the current or emittance is zero.

The next model represents the beam as a drastically elongated 2-dimensional K-V elliptical beam with envelope X_m and Y_m described by the equations

$$X_m'' - \frac{2K}{X_m + Y_m} - \frac{\epsilon_x^2}{X_m^3} = 0, \quad (6)$$

$$Y_m'' - \frac{2K}{X_m + Y_m} - \frac{\epsilon_y^2}{Y_m^3} = 0. \quad (7)$$

Again K is the generalized perveance, but the two emittances of the sheet beam are now described by

$$\epsilon_x = X_0 \sqrt{\frac{2kT_s}{qV_0}} \quad (8)$$

and

$$\epsilon_y = Y_0 \sqrt{\frac{2kT_s}{qV_0}}. \quad (9)$$

The initial minor axis of the ellipse X_0 corresponds to the width of the slit and the initial major axis Y_0 corresponds to the diameter of the full beam. These equations have the advantage of characterizing the beam as having a finite length. On the other hand, the beam is modeled here as an ellipse, not as a sheet.

Lejeune and Aubert³ estimate the finite slit effect by calculating the phase space area that results when a zero emittance beam is measured by slit and ideal detector. The resulting emittance would be

$$\Delta\epsilon = \frac{2R_s 2X_0}{\pi L} \quad (10)$$

In a nonzero emittance beam, this finite slit emittance value would be subtracted from the measured emittance to give a lower limit on the emittance. In our case, this is a very poor estimate of the finite slit effect. We use the sheet beam analysis to better estimate this error. If space charge is equal to zero, the 1-dimensional sheet beam equation has the analytical solution

$$X_m(z) = \left[X_0^2 + 2X_0X'_0 + \left(\frac{\epsilon_x^2}{X_0^2} + X_0'^2 \right) z^2 \right]^{1/2} \quad (11)$$

where X_0 and X'_0 are the initial conditions of the envelope. Since the slit samples upright sections of the phase space of the full beam, the initial slope is zero and the solution at the pinhole position (distance L between slit and pinhole) reduces to

$$X_m(L) = \left[X_0^2 + \frac{3kT_s}{2qV_0} L^2 \right]^{1/2} = gL \left[1 + \left(\frac{X_0}{gL} \right)^2 \right]^{1/2} \quad (12)$$

where g is defined as

$$g = \sqrt{\frac{3kT_s}{2qV_0}}. \quad (13)$$

If X_0 is much less than gL , the solution can be approximated as

$$X_m(L) \approx gL \left[1 + \frac{1}{2} \left(\frac{X_0}{gL} \right)^2 \right]. \quad (14)$$

Thus, the second term in the brackets is the fractional increase in the envelope as a result of the finite width of the slit. For half slit widths of 0.0275 mm, 0.069 mm, 0.102 mm and 0.127 mm, the percent increases in the envelope at the pinhole are respectively 0.04%, 0.3%, 0.6%, and 0.9%. The value for temperature T_s used in this equation is derived from the experimentally extrapolated emittance of 97 mm-mrad.

A more accurate method of predicting the finite slit effect is to perform a superposition of point sources, each with a given velocity distribution, over the slit.⁴ From this, the actual final beam density as a function of its width can be calculated. Assuming a Maxwellian velocity distribution with rms half width of σ at the slit and uniform density n_0 across the slit, the density distribution $n(x)$ at the pinhole plane is given by

$$n(x) = \frac{n_0}{2} \left[\operatorname{erf} \left(\frac{|x| + X_0}{\sqrt{2}\sigma L} \right) - \operatorname{erf} \left(\frac{|x| - X_0}{\sqrt{2}\sigma L} \right) \right] \quad (15)$$

where the error function is defined by

$$\text{erf}(u) = \frac{2}{\sqrt{\pi}} \int_0^u e^{-t^2} dt. \quad (16)$$

Recent work by M. J. Rhee and R. F. Schneider⁵ allows the actual rms value of the velocity distribution to be easily calculated from the measured distribution. Using the extrapolated emittance to derive the value of σ , the percent increases in the rms width of the 0.0275 mm, 0.069 mm, 0.102mm, and 0.127 mm slits are respectively 0.02%, 0.2%, 0.5% and 0.7%. These are slightly lower than the 1-dimensional model and more accurate. The values are so small, however, that the finite slit has an almost negligible effect on the measured emittance.

2. Finite pinhole effect

Another source of geometrical error is caused by the finite size of the pinhole/Faraday cup detector. Lejeune and Aubert estimate this effect by assuming an infinitely thin sheet beamlet of zero emittance is scanned by a pinhole of diameter d_p . The measured phase space area is finite and given approximately by

$$\Delta\epsilon = \frac{2R_s d_p}{\pi L}. \quad (17)$$

Like the finite slit estimate, this value is subtracted from the measured emittance to give a lower limit of emittance. Again, this is not an accurate estimate in this experiment. In our case, the pinhole is small with respect to the expanded sheet beamlets. Furthermore, the profiles being measured have no sharp density gradients. Thus, the pinhole provides an accurate averaging process over a nearly linear range of the distribution being measured. The pinhole effect can be calculated accurately by determining the current passing through the pinhole at a certain location when the beam has a gaussian distribution (very nearly true in this case).

The current detected at each pinhole location s_l is

$$I_s = \int_{s_l - r_p}^{s_l + r_p} \int_{\sqrt{r_p^2 - x^2}}^{\sqrt{r_p^2 + x^2}} n(x, y) dx dy \quad (18)$$

where

$$n(x, y) = n_0 \exp\left(\frac{-x^2}{2\sigma^2}\right) \quad (19)$$

is the sheet beamlet density at the pinhole. This reduces to

$$I_s = \int_{s_l - r_p}^{s_l + r_p} 2n_0 \sqrt{r_p^2 - (s_l - x)^2} \exp\left(\frac{-x^2}{2\sigma^2}\right) dx. \quad (20)$$

By evaluating the integral at many locations of s_l , a density profile can be generated. Fitting this profile to a gaussian reveals an increase in rms width of approximately 0.1% over the range of measured values of gaussian rms width σ . Thus, like the finite slit effect, the finite pinhole effect has negligible effect on measured emittance.

3. Compensation for space charge

Compensation for space charge in the sheet beamlets is accomplished by modeling the beamlet as a 1-dimensional or as a 2-dimensional beam. In order to simulate the sheet beams, it is necessary to make several assumptions. The current density must be derived from the normalized current density profile in Fig. 2. Since all models assume uniform current density, and the profile is measured across the center of the beam, the current density at the beamlet center is used in the simulations. Since the actual beamlets are not uniform, this is a possible source of error. The temperature of the beam used to calculate the simulated sheet beam emittance is estimated in the following way. The temperature of the simulated beamlet is varied until its rms width at the detector plane is equal to the rms width of the measured sheet beamlet. The simulation is again run with the space charge term set to zero. The resultant rms width of the simulated profile at the pinhole is the modified in this way. The results of the modifications are shown in Fig. 4. Note that for both the K-V model and 1-dimensional model the compensation is not sufficient, though the 1-dimensional model seems to be the better of the two. These discrepancies have several possible explanations. First, these models are not exact. They assume uniform density of the sheet beamlets. Though this is true over the widths of most of the beams, the density gradient becomes quite large at the beam edge. In addition, the sheet beamlets are not uniform along their lengths. These models use the current density at the point where the pinhole sweeps across the sheet beam. When the slit is sampling near the center of the full beam, regions of high current density surround the lower current density region at the sheet beam center. As the sheet beam propagates, the regions of higher current density will flow into the regions of lower current density. This would cause an increase in the measured width of the sheet beamlets that would not be present in the simulated beamlets. With regard to the experimental apparatus, perhaps there is some effect that has not been adequately taken into account, such as secondary emission of the slit plate or aperture lensing of the slit, that is causing the beamlets to diverge more than they should. All of these possibilities are being explored.

Figure 5: Schematic of the multiple beam experimental set up.

B. Multiple Beam Experiment Update

Theoretical and numerical studies have shown that nonuniform charge distribution is a major cause of emittance growth in focusing channels. In 1987, we began an experiment to investigate this emittance growth.⁶ The nonuniform charge distribution we chose to study was a five-beamlet configuration formed by masking out a solid round beam. The beam was propagated along the transport channel (consisting of 36 solenoid lenses) where numerous pictures were taken using a CCD camera. The five beamlets were found to merge rather quickly in a distance of about 15 cm, as predicted by theory. Surprisingly, however, at a distance of 101 cm from the aperture plate, an image of the initial distribution was detected on the phosphor screen. No other image was found further down the channel. At the end of the channel, the beam was perfectly round showing no spatial structure, and the emittance was measured using the slit/pinhole method described in the previous section. Figure 5 illustrates the experimental set-up of the Multiple Beam Experiment. This experiment was completed in the spring of 1988. Particle simulation studies done in 1988 confirmed the merging effect, but there was very poor agreement with experiment on the image formation. Further particle simulation studies performed last year produced remarkable agreement with the experimental data. Despite the improved agreement, several inconsistencies between the improved simulation and the experimental data remained and have been the subject of this year's study.

The simulations were done by I. Haber (NRL) and H. Rudd (Berkeley Scholars) using the self-consistent 2-D particle-in-cell code SHIFT-XY.⁷ The initial emittance of the five-beamlet configuration is given by the equation

$$\epsilon_i = R_i \left(\frac{2kT}{qV_0} \right)^{1/2} \quad (21)$$

where

$$R_i = [a^2 - 1.6\delta^2] = 4.67 \text{ mm}. \quad (22)$$

Here, T is the temperature of the beam at the aperture plate. R_i is the initial effective ($2 \times \text{rms}$) radius of the five-beamlets, $a = 1.19 \text{ cm}$ is the beamlet radius, and $\delta = 3a$ is the separation distance between the beamlet centers, as shown in Fig. 5.

The magnetic field produced by a lens was represented by an analytical fit to the measured field profile on axis and Taylor expansion was used to obtain the nonlinear off-axis fields. The analytic on-axis field equation is given by

$$B(z) = B_0 \frac{\exp\left(\frac{-z^2}{2b^2}\right)}{1 + \frac{z^2}{a^2}}. \quad (23)$$

In this equation, a and b are fitting parameters and B_0 is the peak field value. The discrepancies in the simulations center around the value of temperature T in equation (21) and the values of B_0 , a and b in equation (23).

In the original simulation studies that yielded poor agreement, the values of B_0 , a , and b used were 83.2 gauss, 4.4 cm, and 2.29 cm, respectively. The initial temperature at the aperture plate used was the temperature of the cathode. It has been since recognized that the beam undergoes a compression of approximately a factor of two between the cathode and the aperture plate. As a result, there should be a corresponding increase in temperature by a factor of four and of the emittance by a factor of two. The value of emittance used in the original simulation which did not take beam compression into account was 32.4 mm-mrad. By including the beam compression effect, one obtains the more accurate value of 64.8 mm-mrad for the initial emittance of the five-beamlet configuration.

Even with this improved value for emittance, the simulation results remained poor. The location of the image and the beam rotation per period in the simulation did not match the corresponding image location and rotation in the experiment. The beam rotation can be inferred from the phosphor screen pictures, which are especially clear at and near the image location of 101 cm downstream from the anode. When calculating the B_0 necessary to produce the observed rotation, it was found that by increasing the value of B_0 by 10% (i.e. from 83.2 gauss to 91.5 gauss), very close agreement between simulation and experiment occurred.

Using the initial emittance value of 64.8 mm-mrad and $B_0 = 91.5$ gauss yielded the excellent results shown in Fig. 6 and Fig. 7. Figure 6 and Fig. 7 show phosphor screen pictures of the beam and the corresponding shade plots from the simulation at $z = 3.4$ cm, 17.0 cm, 30.7 cm, 44.3 cm, 71.5 cm, and 101 cm. The pictures at 30.7, 44.3, and 71.5 cm are taken at the midpoint between lenses M2 and C1, C1 and C2, and C3 and C4, respectively. Lens M2 is the second matching lens, lenses C1, C2, C3, and C4 are the first through fourth diagonal lenses. The picture at 101 cm is the position of the beam image. The improvements in the simulation results were mainly due to modifications in two of the system parameters - emittance and magnetic field. During this past year, we have attempted to investigate not only the magnetic field discrepancy but also some additional questions concerning the initial emittance.

1. Magnetic field

There are several problems concerning the analytic equation used to model the magnetic fields in the multiple beam experiment. First, each lens is not exactly identical. There is a statistical deviation from the ideal lens model that would average out over many lenses. Unfortunately, the image appears after passing through only seven of these magnets. Thus, it is possible that this model does not accurately predict the electron beam's passage through the channel. Another problem exists in the inaccuracy of the analytic equation itself in modelling the on-axis fields that are greater than about 4 cm from the magnet center. The actual fields die off more slowly than the fitted field. This additional field can potentially cause extra focusing and rotation of the beam. Thus, the fitted field model underestimates the focusing of each solenoid.

At this point in the experiment, disassembly of the entire channel in order to measure each individual solenoid is not possible. Therefore, a study of a solenoid considered to be identical to the ones used in the Multiple Beam Experiment was performed. However, the magnetic field of the solenoid was measured on axis. The measured on-axis magnetic field $B_m(z)$ was found to have a peak field of 82.3 gauss at the magnet current used in the multiple beam experiment. This is even slightly less than the 83.2 gauss assumed for the original experiment and simulation. The new field measurement, therefore, did not provide the 10% increase required to explain the image rotation. Let us now examine the beam rotation found in the experiment.

It was discovered in the multiple beam experiment that the beam rotates $365^\circ \pm 5^\circ$ between $z = 0$ cm (aperture plate) and $z = 101$ cm. The rotation caused by each lens can

EXPERIMENT

SIMULATION

1 cm

1 cm

Figure 6: Phosphor screen pictures and simulation point plots show the separate beamlets at 3.4 cm, the merged beam at 17 cm, and the beam at the channel entrance at 30.7 cm.

Figure 7: Phosphor screen pictures and simulation shade plots show the beam waists at 44.3 cm and 71.5 cm, and the image at 101 cm.

be calculated from the paraxial equation

$$\theta - \theta_0 = \int_{z_0}^z \left(\frac{-qB(z)}{2\gamma m_0 \beta c} + p_\theta m_0 c \beta \gamma r^2 \right) dz \quad (24)$$

where θ_0 is initial angular position, m_0 is mass of an electron, c is the speed of light, and p_θ is the canonical angular momentum. During the multiple beam experiment, the current through the two matching lenses and the channel lenses were measured with the same meter. Thus, it is safe to assume that the ratio of the field peaks with respect to the channel lenses are the same as the ratio of the currents through the lenses. The ratio of matching lens 1 and 2 currents to channel current is 0.989 and 0.891, respectively. The beam travels completely through five channel lenses and then to within 5.8 cm of the center of the sixth channel lens to form the image. Recall that the period of the channel is 13.6 cm. The total rotation of the beam at the image plane is therefore closely calculated by

$$\theta_T = \int_{-L}^L \frac{q(0.989 + 0.891 + 5)B(z)dz}{2\gamma m_0 \beta c} + \int_{-L}^{-5.8} \frac{qB(z)dz}{2\gamma m_0 \beta c} \quad (25)$$

Since we already know the value of θ_T , we can calculate the value of the peak field of $B(z)$ necessary to rotate the beam by an angle of θ_T . By normalizing the magnetic fields, the integrals can be calculated numerically for both the measured field and the analytic field. The value of the peak field B_0 is calculated as the ratio θ_T/θ_N where θ_N is the value of θ_T when the magnetic fields used in the equation are normalized. Using the integration limit $L = 8.9$ cm, where the field has dropped to almost zero value (see Fig. 8a), the value of θ_N for the measured field is 744 radian/Tesla. Thus, the value of B_0 for the measured field would have to be 85.6 gauss in order to generate the measured rotation. Using the same value of L for the analytic field, θ_N is 689 radian/Tesla and B_0 is 92.5 gauss. This is quite close to the value used in the simulations (91.5 gauss). These results show quite clearly the inadequacy of the analytic model. However, since the peak of the test solenoid at the multiple beam operating current is measured as 82.3 gauss while the multiple beam rotation indicates a peak of 85.6 gauss, the magnetic field problem is still not fully explained.

In an effort to find a more accurate equation to represent the measured field, a new fit was tried. Maintaining the same form as in equation (23), a new analytic field $B_2(z)$ was fit to the measured field (called here B_m) by adjusting a and b such that

$$\int_{-L}^L B_m^2(z)dz = \int_{-L}^L B_2^2(z)dz \quad (26)$$

and the peaks were assumed to be identical (i.e. 82.3 gauss). This method of fitting was used because the average focusing strength κ of a solenoid is proportional to the integral of the square of the magnetic field. The values of a and b were found to be 44.24 mm and 25.14 mm

Figure 8: a) Plot of on-axis magnetic field for the measured test magnet field B_m , fitted field B_2 , and original analytic fit B_1 . b) Plot of axial field 9.7 mm off-axis for the three fields in part a).

respectively. The measured field, the new fitted field, and the old fitted field (called $B_1(z)$) are shown in Fig. 8a. The ratio of the focusing strengths is 1.04. This implies that, in order for the focusing strengths to be equal, the amplitude of $B_1(z)$ would have to be increased by 2%. Off-axis axial fields generated by $B_1(z)$, $B_2(z)$, and the corresponding measured field are shown in Fig. 8b.

Table 1 summarizes the comparison between the different magnetic fields under discussion: the measured field B_m , the original field B_1 that resulted in poor agreement, the field B_{sim} used in the simulation (see Fig. 6 and Fig. 7) that resulted in good experimental agreement, and the new fitted field B_2 . Except for B_{sim} , the values of B_0 shown are the ones that correspond to the multiple beam experimental conditions. In the case of B_{sim} , B_0 is simply 10% higher than the original field B_1 . The zero current phase advance σ_0 was calculated using the equation.

$$\sigma_0 = \int_{-s/2}^{s/2} \frac{ds}{w^2} \quad (27)$$

where w is defined by

$$w'' + \left(\frac{qB(z)}{2m_0c\beta\gamma} \right)^2 w - \frac{1}{w^3} = 0 \quad (28)$$

Initial conditions w_0 and w'_0 were determined by satisfying matched beam conditions and $B(z)$ was determined by equation (23). It is defined here for later reference. The averages of the square of the fields (proportional to the focusing strength κ) defined by

$$(B^2)_{avg} = \frac{1}{L} \int_{-L/2}^{L/2} B^2(z) dz \quad (29)$$

are calculated with $L = 17.8$ cm for the on-axis fields and 12.0 cm for the off-axis fields. Also shown are the beam rotations through one lens caused by each field, and the total beam rotation through the channel to the image plane. Notice that there is substantial error in the focusing strengths of the analytically generated fields. The rotation of the beamlets over a single channel period can be measured through the sixth channel lens ($z = 98.3$ cm to $z = 112.3$ cm) to be $52^\circ \pm 4^\circ$. Unfortunately, the resolution in angle is not fine enough to differentiate the different cases. With respect to the beam rotation, the field used in the simulation agrees well with results of the multiple beam data but the agreement breaks down when comparing the average squared field with that of the measured test solenoid. The new measurements show improvement over the original field B_1 , but substantial discrepancy still remains. Despite this, the matching of the beam rotation gives strong justification to the 10% increase in the field that yielded poor results and it is this result that can be most trusted. The presence of some gross error in the channel solenoid currents or some other parameter, though not likely, is always a possibility.

Table 1: Summary of the comparison between the different magnetic fields under discussion: the measured field B_m , the original field B_1 that resulted in poor agreement, the field B_{sim} used in the simulation (see Fig. 6 and Fig. 7) that resulted in good experimental agreement, and the new fitted field B_2 . Except for B_{sim} , the values of B_0 shown are the ones that correspond to the multiple beam experimental conditions. In the case of B_{sim} , B_0 is simply 10% higher than the original field B_1 . Shown are values of σ_0 and average square of the magnetic field on and off-axis. Also shown are the beam rotations through one lens caused by each field, and the total beam rotation at the image plane of each.

	B_m	B_1	B_{sim}	B_2
B_0 [gauss]	82.3	83.2	91.5	82.3
σ_0	—	68.9°	76.3°	70.4°
$(B^2)_{avg}$ [gauss ²] ($r = 0$ mm)	1347	1289	1560	1347
$(B^2)_{avg}$ [gauss ²] ($r = 9.7$ mm)	2124	1962	2374	2079
θ_{lens}	50.0°	47.7°	52.5°	50.0°
θ_{image}	344°	328°	361°	344°

2. Initial Emittance

Though good experimental agreement was attained by increasing the magnetic field, another inconsistency exists in the simulations and involves the initial emittance. When the simulations were first done, the radius and emittance of the beam at the aperture plate had not been accurately measured. Thus, the intrinsic emittance was used and a compression of exactly two between cathode and anode was assumed. This initial emittance value was 64.8 mm-mrad and, when used with the 10% increase in magnetic field, yielded excellent results. With the completion of the test stand measurements, fairly reliable values for both the initial beam radius and temperature were established. Thus, a more accurate value of emittance for the five-beamlet distribution can be derived.

The test stand measurements discussed earlier determined the emittance of the full beam to be 97 mm-mrad 45 mm from the aperture plate. To apply this to the multiple beam experiment, we want to deduce the temperature of the beam at the aperture plate. To do so, first calculate the average temperature of the beam at 45 mm. The effective radius of the beam is given by

$$R = \sqrt{2} R_{\text{rms}} \quad (30)$$

R_{rms} can be calculated numerically from the beam profile (see Fig. 2) to be 4.85 mm. This corresponds to $R_s = 6.86$ mm. Assuming the emittance of the beam is accurately described by

$$\epsilon = R \sqrt{\frac{2kT}{qV_0}}, \quad (31)$$

the average temperature of the beam can be calculated as 0.496 eV. The beam radius at the aperture plate was measured last year to be 6.1 mm. This corresponds to an effective radius R_a of about 6.0 mm. Assuming there is no emittance growth between the aperture plate and the slit, the average temperature at the aperture plate kT_a can be calculated as

$$kT_a = kT_s \left(\frac{R_s}{R_a} \right)^2 \quad (32)$$

where T_s is the beam temperature at the slit. This yields a value of 0.648 eV. The measured temperature of the cathode was 0.115 eV. This corresponds to an emittance growth of 1.12, like the emittance measurement. The cathode temperature in the multiple beam experiment was 0.122 eV. Assuming the beam radius at the multiple beam aperture is the same as in the test stand (good assumption since the beam is highly space charge dominated) and the emittance growth is identical to that found in the test stand, the effective emittance of the five-beamlet distribution can be calculated to be about 76 mm-mrad.

Such a change in emittance has a substantial effect on the location of the image. The theory of emittance growth due to space charge homogenization^{8,9} predicts an emittance

growth of

$$\frac{\epsilon_f}{\epsilon_i} = \left[1 + \frac{U}{2w_0} \left(\frac{\sigma_0^2}{\sigma^2} - 1 \right) \right]^{1/2} \quad (33)$$

or the alternate form

$$\frac{\epsilon_f}{\epsilon_i} = \left[1 + \frac{KR^2U}{2\epsilon_i^2w_0} \right]^{1/2} \quad (34)$$

where σ_0 is the phase advance of particle trajectories per period with no space charge, σ is the phase advance per period with space charge, and U/w_0 is a dimensionless quantity depending only on the shape of the nonuniform distribution. In this experiment, $U/w_0 = 0.2656$. For the simulation that yielded close experimental agreement (i.e. 10% increase in B_0 and $\epsilon_i = 64.8$ mm-mrad) the emittance growth was 1.55. This corresponds to a final emittance of 100 mm-mrad. The measured emittance at the end of the channel (520 cm downstream of aperture plate) is $108 \text{ mm-mrad} \pm \text{mm-mrad}$. With regard to the final emittance, the simulation and experiment agree quite reasonably in view of the extra emittance growth expected from nonlinearities of the solenoids and beam offset.

One interesting fact arises from the formation of the image of the initial distribution. From simple particle optics, an image is formed at the half betatron wavelength of the beam particles. An equation for the image locations z_i as a function of period length S and phase advance per period with space charge σ is easily derived.

$$z_i = \frac{n\pi S}{\sigma} \quad (35)$$

Here $n = 1$ is the first image, $n = 2$ is the second image and so on. Since the image position is known from the experiment to be 101 cm, the value of σ can be easily calculated as 24.2° . This value is true for both experiment and simulation. From smooth approximation theory¹⁰, a value for emittance can be calculated from the equation

$$\epsilon = \frac{KS}{\sigma_0} \left(\frac{\sigma}{\sigma_0} \right) \frac{1}{[1 - (\sigma/\sigma_0)^2]}. \quad (36)$$

For the simulation which yielded the good results, with $\sigma = 24.2^\circ$, $K = 1.877 \times 10^{-3}$, $S = 0.136$ m, and $\sigma_0 = 76.3^\circ$, the value of emittance turns out to be 67.6 mm-mrad, which is very close to the initial emittance of 64.8 mm-mrad. This seems to indicate that the particles that form the image have not yet experienced emittance growth due to the homogenization of space charge. In a parameter range where smooth approximation theory is accurate, this provides an accurate method of determining the initial emittance. Note that the second image does not form. The simulation clearly predicts that no additional emittance growth occurs after that predicted by space charge homogenization theory. This growth occurs in approximately a quarter of a plasma wavelength. In this experiment, the emittance growth

occurs within 17 cm of the aperture. The presence of a first image and absence of any others indicates a longer time scale for further spatial homogenization that is not accompanied by emittance growth.

Since there is some doubt as to the veracity of the experimental value of B_0 , it is interesting to see the effect of both focusing strength (σ_0) and emittance on image location using the smooth approximation theory. The phase advance with space charge is given by the equation

$$\sigma = \sigma_0 (\sqrt{u^2 + 1} - u). \quad (37)$$

where

$$u = \frac{KS}{2\sigma_0\epsilon_i} \quad (38)$$

It is obvious that, with regard to image position, the two variables behave very similarly. Increase in either σ_0 or ϵ_i causes the image to move closer to the aperture plate. Thus, if the image position is to be held constant and the focusing strength is increased, the emittance must decrease and vice versa. If the parameters used in the simulation are input into these formulas, the predicted image location is 105 cm. Since a sharp image appears from about 98 cm to 112 cm, this is quite an acceptable value. Assuming the initial emittance derived from the test stand (76 mm-mrad) and the value of σ_0 that results in acceptable beam rotation (76.3°), the image location is predicted as 92 cm. This is not seen experimentally.

Thus, a dilemma has been reached which at this point has not been resolved. Rotation of the beamlet image dictates a certain value of the magnetic field (10% increase over previous values). According to experimental measurements, on the other hand, the initial emittance used in the simulation appears to be too low by about 12%. Smooth approximation theory predicts that only one of these values can be increased and still maintain the correct image location.

More extensive multiple beam experiments are planned for next year. Hopefully some of these discrepancies can be solved. In addition, future work is planned to study the regime where estimation of the initial emittance via image formation remains accurate.

C. Longitudinal Pulse Compression and Instability Experiment

Considerable progress has been made in the effort to improve the performance characteristics of the new electron beam injector and its components. Notable achievements over the past year are as follows:

1. Preliminary test of the new electron beam injector

The new electron beam injector, which consists of the variable-perveance gridded electron gun, three matching lenses, and the induction acceleration module, was assembled and preliminarily tested last spring. The results were reported at the American Physical Society Conference, April 16-19, 1990, in Washington, DC.

2. Electron gun improvements

The design of the variable-perveance gridded electron gun and its general performance characteristics were described in detail in last year's progress report. Since the cathode assembly ML-EE55 used for the gun was no longer available and had to be replaced by a new product, substantial changes and performance improvements were made on the gun. Here is a list of the main points.

1. The old ML-EE55 oxide cathode was replaced by the Y646B dispenser cathode. The Y646B cathode assembly also has a planar triode configuration consisting of the heater, cathode, and grid. The cathode is an indirectly heated disc of area 0.5 cm^2 , smaller than that of ML-EE55. The Y-646B has 40×40 woven wire mesh grid with spacing 150 microns at a wire size of 1 mil diameter.
2. The distance between the grid and the focusing electrode was reduced from the original 1 mm to about 0.13 mm. This has dramatically improved, without worsening the problem of the thermal expansion and conductivity in the structure, the field uniformity in the region before the grid and has increased the beam current controlling sensitivity of the cathode-grid pulse.
3. The anode was shaped to conform to the ideal Pierce geometry at the A-K gap of four times the cathode radius. This is the position close to the operating condition for the pulse compression experiment. This position should theoretically produce parallel flow and result in the minimum emittance.
4. The A-K gap was modified to range from 0.93 cm to 2.3 cm, which results in the gun perveance ranging from 0.22×10^{-6} to $1.35 \times 10^{-6} \text{ A V}^{-3/2}$. It is common practice to regard a beam as space charge dominated if its perveance is above 1 microperv. This is the case in the improved gun performance.
5. The original differential Rogowski coil was replaced by an integrating one. This allowed direct measurement of the beam current without the external amplifier and integrator.

Figure 9: Beam current versus anode voltage where the smooth curve is from Child's law.

A simple filter was employed to smooth the output signal of the coil. Though this slowed down the rise time of the signal pulse, the amplitude of the coil output signal has a linear relationship with the magnitude of the beam current.

Figure 9 shows the current as a function of the anode voltage, and Fig. 10 plots the beam current as a function of the A-K gap. In comparison with the previous results, much higher current and better agreement with the theory was achieved in the improved gun performance.

Besides the improvements of the hardware of the gun, theoretical study on the beam emittance measurement by pepperpot method was also made. The result is reported in a paper entitled "Beam Emittance Measurement by Pepper-pot Method." which will be submitted for publication

The emittance measurement for the beam is underway and the new results from this analysis will be reported elsewhere.

3. Induction linac testing

The induction acceleration module testing was completed last summer. The test results were presented at the 1990 Linear Accelerator Conference, September 10-14, 1990, Albuquerque, NM. A paper entitled "Design and Performance Characteristics of a Compact Induction Acceleration Module for Longitudinal Pulse Compression Experiments," which details all the aspects of this device, has been accepted for publication in *Nuclear Instruments and Methods in Physical Research, Section A*

Figure 10: Beam current versus A-K gap where the current is measured at a fixed anode voltage and the smooth curve is generated from Child's law.

4. Beam matching studies with computer code

A modified K-V code has been used to simulate the beam transport property for the compression experiment, which takes the induction linac and all matching lenses into account. Since the beam has an energy (i.e., velocity) spread while focusing conditions are fixed, it is actually impossible to match the whole beam. However, there exist some operation conditions at which the whole beam, varying with the energy from 2.5 to 7.5 keV, is mismatched, but still can be confined through the channel. The simulation results are shown in Fig. 11.

5. Resistive-wall instability study

The resistive wall instability concerns beam physics in many applications such as heavy inertial fusion by induction accelerators. We have found we are in a unique position to study this phenomena experimentally using our electron beam injector with some associated facilities. Over the last few months the effort was made to understand theoretically the mechanism of the instability, to design the experiment, and to produce a beam transport pipe with resistive-wall. A paper entitled "Experimental Study of the Longitudinal Instability for Beam Transport," which describes this resistive-wall instability, was presented at the International Symposium on Heavy Ion Inertial Fusion, December 3-6, 1990, in Monterey, CA.

Figure 11: The beam envelope along the injector and periodic channel under certain focusing conditions, but at different induction linac gap voltages, where (a) is for electrons with 2.5 keV energy at the front of the beam (no gap voltage applied, $V_g = 0$), (b) is for ones with 3.5 keV ($V_g = 1$ kV), and (c) is for ones with 7.5 keV at the tail of the beam ($V_g = 5$ kV).

D. Theoretical Studies

The advanced accelerator applications, such as linear colliders, free electron lasers, heavy ion inertial fusion, and high brightness H^- beams for the SSC or defense, require high-brightness beams that are pushing the state of the art. One of the key issues is to obtain an understanding and control of all effects that cause emittance growth. Our research program at the University of Maryland has played the leading role in identifying space charge nonuniformity as a major culprit for the significant emittance growth observed in experiments and accelerators with high brightness beams. Nevertheless, several other effects, notably beam mismatch and off-centering, known to also cause emittance growth, had so far not been explained by a self-consistent theory of high-brightness beams. In what can be considered a major breakthrough, we have succeeded during the past year to close this gap.

The crucial starting point for this new theoretical model is the recognition that the stationary (equilibrium) beam in a focusing channel or accelerator is characterized by two conditions:

1. There is a balance between the average (rms) applied focusing forces and the repulsive rms forces due to space charge and emittance.
2. The total transverse energy (transverse kinetic particle energy + potential energy due to the applied force + space-charge field energy) is a minimum.

Any difference from this stationary state, such as mismatch of the density profile, mismatch in rms radius or transverse rms velocity and off-centering, increases the total transverse energy of the beam. The additional energy ΔE constitutes "free energy" which can cause emittance growth. Our new model provides the relatively simple mathematical relations for calculating ΔE and for determining the possible emittance growth resulting from ΔE . The key feature in our model is that we do not attempt to solve the highly nonlinear equations describing the internal particle dynamics of a nonstationary space-charge dominated beam. (To a large extent this attempt has so far been unsuccessful and beam dynamicists rely almost exclusively on particle simulation to study the physics. From such simulation, however, it is known that nonstationary beams tend to evolve towards a final stationary state with increased emittances.) Instead we use average force balance and energy conservation to obtain the phase space boundaries within which the particle distribution must be confined.

Our model provides the analytical framework that explains this effect and, moreover, gives an accurate estimate for the emittance growth that occurs as the beam converges towards the final stationary state. The new theory is presented in a paper entitled "Free Energy and Emittance Growth in Nonstationary Charged Particle Beams," which was submitted for publication to *Physical Review Letters*

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TASK B

**Study of Collective Ion Acceleration by Intense Electron Beams
and Pseudospark Produced High Brightness Electron Beams**

I. Introduction and Synopsis

Task B of this contract was concerned with experimental and theoretical studies of intense pulsed-power electron beam systems and their application to advanced accelerator concepts. In particular, work supported under Task B focused around two experimental projects; the laser-controlled collective ion accelerator and studies of high-brightness electron beams produced in pseudospark discharges. Although progress on the laser-controlled collective accelerator was steady and encouraging to date (see section II.A), our quest for even higher accelerated ion energies was hampered by the voltage and pulse duration capabilities of our existing intense beam sources. Based on experiments conducted on pseudospark devices in our laboratory over the past two years, we believe that this new class of high-brightness electron beam source may have important applications in such diverse areas as linear colliders, free electron lasers, rf sources, high power switching systems, and materials processing. In fact, our experiments indicate that high-brightness beams can be generated with current densities of 10^3 to 10^6 amps/cm², orders of magnitude higher than competing low-emittance electron sources. A more complete introduction to this area of research is included below.

A. Laser Controlled Collective Ion Accelerator Experiments

In these experiments, control over the motion of a virtual cathode at the front of an intense relativistic electron beam was achieved by an ionization channel generated by the time sequenced formation of plasma clouds along the drift tube. This time-sequenced ionization channel allowed for beam neutralization and propagation in a controlled manner. Extensive experimental work was completed on a 50 cm long system in which beamfront control and ion acceleration was demonstrated at an effective accelerating gradient of 40 MV/m. Attempts to accelerate ions at a higher accelerating gradient of 90 MV/m in this system were unsuccessful. Numerical simulations indicated that we are unable to maintain a sufficiently high electric field at the moving virtual cathode to accelerate with this gradient. Over the past year, a 100 cm experiment was designed with the aid of our numerical simulation code, constructed, optical system tests completed, and beamfront control confirmed for two different accelerating gradients. These studies are detailed in section II.A.

Additional theoretical studies verified the existence of an equilibria between co-moving electron and ion beams. The Bennett profile of each species was examined with respect to the ion hose instability.

B. Pseudospark Discharge Experiments

During the last decade, considerable research has been conducted on pseudospark discharges of the type first explored by Christiansen and Schultheiss in 1978.¹ Interest in this novel discharge configuration has been driven by potential applications of such discharges to such areas as high power switch development. In addition, the observation of high current density electron beams generated by such discharges has spurred interest in their possible application as high-brightness electron beam sources for such diverse uses as electron beam lithography and plasma processing, for example. Moreover, the development of novel coherent radiation sources such as the Free Electron Laser (FEL) and the challenges faced by designers of next-generation linear collider have spurred renewed interest in new methods for producing low-emittance, high brightness electron beams. Both FELs and electron-positron colliders, for example, require electron beams of exceptionally low emittance, high brightness, and low energy spread. Other applications, such as high power microwave tubes would also clearly benefit from the development of new high-brightness beam sources, particularly if total beam current and energy could approach that currently achieved in high power pulse line accelerators.

To date, most electron beams for such applications are produced either by thermionic diodes, where beam quality can be kept reasonably high but current density is usually limited to values below 20 A/cm², or by field emission diodes where current density can be orders of magnitude higher but high beam quality is difficult to achieve. During the last decade, however, electron beams with high current density have been generated in pseudospark discharges of the type first explored by Christiansen and Schultheiss. In an experiment in our laboratory, for example, we measured the emittance of such a beam and found that the associated brightness was extremely high, as discussed in section II.B. If our results and the beam qualities already demonstrated in other low voltage pseudospark devices can be maintained as they are scaled to higher beam energies and currents, it is clear that pseudospark discharges may be the key to the production of high power, high-brightness electron beams for a wide variety of potential applications.

Pseudospark discharges are unusual in that they operate on the low pressure side of the Paschen breakdown curve minimum, as shown in Fig. 1. Since breakdown in any device exhibiting Paschen behavior is a function of the product of ambient pressure and electrode spacing, operating on the low pd side of the Paschen minimum means that the breakdown voltage is actually lower for long electrode gaps than for short ones. Consequently, pseudospark devices are designed to initiate breakdown along the longest path length available between the two electrodes, as shown in Fig. 2. Although they were first studied as potential

Figure 1: Typical Paschen curve

electron and ion beam sources, pseudospark devices have for the most part been investigated as attractive alternatives to present high current switches such as thyrotrons, and indeed their performance in this area has been extremely promising.

1. Historical background and principle

The "pseudospark" phenomenon was first reported by Christiansen and Schultheiss¹ in 1978 as a fast low pressure gas discharge which occurs in a special device, called a "pseudospark chamber". The breakdown time of this discharge was found to be similar to that of a high pressure spark gap although the two mechanisms appear to be totally different.^{2,3} A single-gap pseudospark chamber consists of an anode with a center hole and a hollow-cathode, with both electrodes separated by a thin insulating washer, as shown in Fig. 2.

The pseudospark chamber is perhaps best described as a combination of a restricted linear discharge geometry and a hollow cathode. The breakdown time depends on the operating point on the breakdown curve, i.e., fast breakdown occurs at higher voltage (lower pressure). In typical operation a characteristic breakdown curve similar to the Paschen curve for parallel electrodes is obtained. The pseudospark is initiated on the left side of this breakdown curve. In this regime the gas discharge occurs along the longest possible path which, in the case of a pseudospark geometry, is found on the system axis. The voltage holdoff capability of one gap is limited to 40 to 50 kV due to the onset of surface flashover and field emission.

Figure 2: Basic scheme of a single gap pseudospark chamber

This holdoff voltage, however, can be enhanced by adding additional stacks of intermediate electrode and insulator discs forming a multigap chamber, as shown in Fig. 3.

Although the pseudospark phenomenon was discovered more than a decade ago, its mechanisms are still not fully understood. Some of the most interesting information has been obtained through spectroscopic observations.^{2,3} The discharge starts with a high impedance Townsend predischage, which lasts for a few microseconds before the main breakdown occurs. A light emitting zone has been observed, starting from the anode and moving toward the cathode with a constant velocity of about $1 \text{ cm}/\mu\text{s}$ measured in 150 mTorr of argon gas. The current rises from 10 A to a few hundred A until a stationary high voltage glow discharge with constant current is achieved. At this point the light emitting zone becomes stationary at a distance of a few mm in front of the cathode. During the predischage, a high reduced electrical field strength E/n (where E is the electric field and n the particle density) of over 10^{-14} V cm^2 is produced resulting in high energy runaway electrons. This explains the electron emission through the anode hole during the predischage. At the same time positive ions drift toward the cathode creating a positive space charge region.

When the positive space charge on axis at the cathode hole reaches a critical value, the hollow cathode discharge is ignited and the main discharge takes place. Ionization waves moving into the hollow cathode at velocities of approximately 10^6 m/s are observed. The light intensity in the gap between the anode and cathode rises and an intense electron beam is ejected at the anode. When the voltage drops, within about 10 ns, the current rapidly rises at a rate of up to $8 \times 10^{11} \text{ A/s}$ and current densities exceeding 10^5 A/cm^2 are observed.⁴ These extremely high current densities in a glow discharge are inconsistent with a cold cathode emission process, and could possibly be attributed to field enhanced emission from the melted electrode surface.⁵ The first numerical simulation results of the pre- and early main discharge phases of the pseudospark were reported⁶ recently.

It is possible to trigger the main discharge in the hollow cathode region with subnanosecond jitter, without the long predischage, when a sufficient amount of charge carriers is

Figure 3: Multigap Pseudospark Chamber

provided in the hollow cathode. This can be achieved either by initiating a surface flash over,⁷ by providing a pulsed glow discharge⁸ or illumination with UV-light.⁹

The pseudospark discharge was initially introduced as a novel type of charged particle beam source for electron and ion beams,¹ and pinched electron beams with current densities up to 10^6 A/cm² and power densities up to 10^9 W/cm² have been observed.⁷ In addition to applications mentioned previously, such high-power electron beams can be used for material processing (for example drilling holes in metal targets or evaporation of semiconductors and isolators to produce layers of high temperature superconducting materials¹⁰). With the emission of the intense electron beam, pulsed microwaves¹¹, and X-rays are also observed.^{2,12}

2. Related work at other labs

The pseudospark is easily triggered with high precision in the hollow cathode region. This behavior favored the development of high power pseudospark switches which were first studied at CERN.⁷ Since then, pseudospark switches for various applications have been developed. Current pulses up to 200 kA with a jitter of ~ 100 ns have been switched at 0.3 Hz for over 500,000 shots without significant electrode deterioration.¹³ With a pseudospark switch designed for laser circuits a jitter of 4 ns at 25 kA peak current and 100 Hz repetition rate was measured.⁴ Multichannel pseudospark switches¹⁴ have achieved a current rise time rate of 2.4×10^{12} A/s with a peak current of 15 kA and a jitter of 1 ns.

Work on the pseudospark phenomenon in the US is mainly performed by Prof. M. Gundersen and his group at USC, where the application of the pseudospark discharge in a thyatron-type switch triggered by UV light is currently under study (this is the Back-Lighted Thyatron or BLT).¹⁵ Additional work is also underway at Old Dominion University.

3. Potential applications for high-brightness electron beams and discussion of competing sources

Advanced accelerator applications such as e^+e^- linear colliders or high-power Free Electron Lasers (FELs) require electron beams whose brightness must be considerably higher than what has been achieved in existing accelerators or experiments. For optimum FEL operation, the required beam emittance scales linearly with the wavelength: the shorter the wavelength, the smaller must be the emittance. These requirements motivate ongoing research on the physics of high-brightness sources and on the development of new sources with better performance characteristics. "Source" will be defined here as the "injector," i.e., electron gun plus beam manipulation system, that produces an electron beam of the desired characteristics (current, pulse length, emittance, etc.) for the main accelerator (usually an rf or induction

linac). The electron emitter used exclusively so far is the thermionic cathode. It has an intrinsic current density limit of about $J_c = 20 \text{ A/cm}^2$. This, in turn, also limits the intrinsic brightness $\mathcal{B}$ of the electron beam, since $\mathcal{B} \propto J_c$.

As an example of a high-brightness accelerator let us consider the $50 \text{ GeV} \times 50 \text{ GeV}$ linear collider (SLC)¹⁶ being developed at SLAC. It requires electron bunches of 5×10^{10} particles with bunch length of about 16 picoseconds (peak currents of about 500 amperes) and normalized emittance of $3 \times 10^{-4} \text{ m-rad}$ from the injector. To get the low emittance, a thermionic cathode with small radius is used which produces an electron beam current of a few amperes, with pulse length of about 2 nanoseconds. This beam is then compressed by a factor > 100 in a subharmonic bunching system to achieve the desired pulse length for the linac. Emittance increase during pulse compression has been a major challenge for this design. The rf linac being developed at Boeing¹⁷ to drive an FEL has beam requirements that are not much different from those for the SLC. It, too, uses a thermionic cathode and a subharmonic buncher strategy similar to the SLC injector. In both the SLC and the Boeing injector the brightness is increased by the bunching process which raises the peak current by the compression factor. However, in practice the emittance increase during subharmonic bunching, partly offsets this gain.

A different approach is being pursued at Livermore where the induction linac capable of handling high peak currents is used as a driver for an FEL¹⁸ and for various other experiments such as the two-beam accelerator¹⁹ and the relativistic klystron²⁰ for future linear colliders. There, large-area thermionic cathodes in diodes with accelerating voltages in the megavolt range produce electron currents of several kiloamperes with pulse lengths of about 50 ns. While the emittance is larger than in the case of the small-area cathodes of rf linac injectors, the intrinsic brightness is the same since J_c is the same. (However, the net brightness gain in subharmonic bunching gives the rf linac injector an advantage over the induction linac.)

The intrinsic brightness limitations of thermionic cathodes as well as the emittance degradation observed in subharmonic bunchers have motivated a search for other sources that do not suffer from these shortcomings. A new concept, the laser photo-cathode, also known as the "rf gun", pioneered at Los Alamos²¹ and also being developed at Brookhaven²² and other places, appears capable of operating at current densities of $J_c \approx 600 \text{ A/cm}^2$ and hence greatly enhanced intrinsic brightness. The research at Brookhaven, combining the photocathode with high-voltage, pulsed-power switching technology, aims at even higher values of current density ($J_c > 1000 \text{ A/cm}^2$) and brightness.

One major technological problem with the photocathode appears to be the lifetime of the cathodes. Further development is needed to determine whether this problem can be solved. Given the limitations of thermionic cathodes and the fact that the photocathode is

still in a developmental stage, there is considerable interest in other possibilities and new ideas of generating high-brightness electron beams. In this regard, the pseudospark device is a very attractive source for such beams. With a proper geometry, this device can produce high-current electron beams with very small diameters. Peak current densities as high as 10^5 to 10^6 A/cm² have been inferred from measurements of the holes drilled with these beams through metal foils.⁷ Recently, the first measurements of the emittance and brightness of a pseudospark electron beam were carried out in our Charged Particle Beam Laboratory at the University of Maryland.²³ Though the beam had only modest values of voltage and current – 25 kV and 100 A, respectively – the experimental results are very impressive. They indicate a brightness that is better by almost an order of magnitude than what has been achieved in beams from thermionic cathodes so far. To put these results in perspective, we shall briefly discuss the theoretical concepts of emittance and brightness of a beam.

The intrinsic emittance of the electron beam emitted by a thermionic cathode is determined by the cathode radius, r_c , and the transverse velocity spread which, in turn, depends on the cathode temperature, kT . Since the electrons in the cathode have a Maxwellian velocity distribution, $f(v) = f_0 \exp(-mv^2/2kT)$, they emerge with an rms velocity of $\bar{v}_x = \sqrt{kT/m}$, where m is the rest mass. Assuming uniform current density across the cathode, the beam has an rms width of $\bar{x} = r_c/2$. The “normalized” rms emittance is then simply given by the product of $\bar{x}\bar{\beta}_x$, where $\bar{\beta}_x = \bar{v}_x/c$, i.e., $\bar{\epsilon}_n = 0.5r_c(kT/mc^2)^{1/2}$. Following Lapostolle’s suggestion, we prefer to use the effective normalized emittance defined as $\epsilon_n = 4\bar{\epsilon}_n = 4[\langle x^2 \rangle \langle \beta_x^2 \rangle - \langle x\beta_x \rangle^2]^{1/2}$. At the cathode the term $\langle x\beta_x \rangle^2$ is zero and hence,

$$\epsilon_n = 2r_c(kT/mc^2)^{1/2}[\text{m} - \text{rad}]. \quad (1)$$

This definition has the advantage that for a uniform beam it is identical with the total emittance while in the case of nonuniform beams it comprises typically between 90% to 100% of the particles. Note that we define the area of the phase-space ellipse as $A = \epsilon_n \pi$, i.e., the “emittance” ϵ_n does not contain the factor π .

The intrinsic normalized brightness of a round beam is defined as

$$B_n = 2I/\epsilon_n^2 \pi^2 [A/(\text{m} - \text{rad})^2]. \quad (2)$$

In an ideal system with linear forces and no longitudinal compression, the current, I , and the normalized emittance, ϵ_n , remain constant; hence the brightness, B_n , is also constant. Eqs. (1) and (2) thus represent lower and upper theoretical limits, respectively, that can serve as figures of merits with which to compare the actual values achieved in practice. Since

current density at the cathode, J_c , is limited by technological constraints, we use $I = J_c r_c^2 \pi$ in (2) and substitute for r_c^2 from (1) to get

$$B_n = (J_c/2\pi)(mc^2/kT) \quad (3)$$

For thermionic cathodes one has a maximum current density of about $J_c = 20 \text{ A/cm}^2$ and a typical temperature of $kT = 0.1 \text{ eV}$. Consequently, one obtains for this case from (3) an upper brightness limit of

$$B_n = 1.63 \times 10^{11} \text{ A/(m - rad)}^2. \quad (4)$$

Note that this theoretical limit is independent of beam current and cathode radius. A cathode with larger radius, r_c , produces a higher current ($I \sim r_c^2$), but the normalized brightness remains the same. However, the emittance, ε_n , increases linearly with r_c . If the emittance is given for a particular application, the radius, r_c , and hence the beam current, I , are fixed. Subharmonic bunching must then be used to increase the current to achieve the desired pulse length. Alternately, one can replace the thermionic cathode by a laser photocathode, as discussed above. In photocathodes, the average kinetic energy of the emitted electrons is defined by the difference of photon energy, $h\nu$, and workfunction, W , and is higher than the temperature of the thermionic cathode depending on the type of material used. With semiconductors one might get $kT = h\nu - W = 1 \text{ eV}$, for metals several eV. On the other hand, the current density may be considerably higher than for thermionic cathodes. Using $J_c = 600 \text{ A/cm}^2$, $kT = 1 \text{ eV}$, one obtains an intrinsic brightness for photocathodes of

$$B_n = 4.86 \times 10^{11} \text{ A/(m - rad)}^2. \quad (5)$$

This is a factor of 3 greater than the brightness of a thermionic beam. More important, however, is the fact that the emittance of the beam from such a photocathode is $\sqrt{600/20} \approx 5.5$ times smaller than for a thermionic beam with the same current. The advantage of the photocathode is apparent from these numbers. However, it must be pointed out that photocathodes are under development. Many problems, such as life time, operation at high average power, etc. remain to be solved before it will become a practical device.

The actual brightness values observed in experiments or numerical simulations are significantly below the intrinsic limits given above – both for thermionic cathodes as well as for photocathodes. These brightness limitations are caused for the most part by nonlinear external forces²⁴, nonlinear space charge forces, and time-varying forces (e.g., rf phase variation) as discussed in Ref. 25.

One finds that the actual brightness values measured in photocathode experiments are in the range of 10^9 to 10^{10} A/(m-rad)², i.e., one to two orders of magnitude below the intrinsic values of thermionic and photocathodes.²⁵ The measurements with our modest pseudospark beam²³ described in the next section show a brightness of about 4×10^{10} A/(m-rad)² at a current density of 250 A/cm². Since one can operate at much higher current densities at higher voltages, one would expect considerable improvements in the future research being proposed here. In fact, a major research goal would be to determine the scaling of current density, emittance, and beam brightness with voltage, geometry, gas composition, and pressure of the pseudospark device. This involves obtaining an understanding of the discharge and beam physics in the pseudospark. A second important goal would be to extract the electron beam from the gas drift tube into vacuum and study how to focus it before space-charge repulsion would blow it up. It is obvious that electron beams with such high current density and brightness must be quickly accelerated to high energy while being focused effectively at the same time. A scheme in which the electron beam from the pseudospark chamber is injected directly into a high-gradient rf cavity, similar to the photocathode experiment at Los Alamos,²¹ looks very attractive and would be examined.

4. Summary of recent pseudospark experiments at Maryland

During the past contract period, experimental efforts in this area have been increased and a second, higher voltage pseudospark experiment has been designed, constructed, and tested. This device was designed to bridge the gap between the lower voltage experiments previously conducted in our laboratory and planned experiments on our high power pulse line accelerators. These experiments, described in detail in section II.C, indicate that a pseudospark discharge with a charging voltage of 50 kV and resultant electron beam currents in excess of 1 kA can be reliably operated over a range of experimental conditions. Measurements of electron beam emittance and brightness are currently in progress.

II. Research Progress

A. Experimental Research on Collective Field Accelerators

1. Laser controlled beamfront accelerator experiments

The basic concept behind the Laser Controlled Beamfront Experiment is shown in Fig. 4. An intense relativistic electron beam is injected into an evacuated drift tube at a current level several times the vacuum space charge limit, given approximately by

$$I_t = \frac{17,000(\gamma_o^{2/3} - 1)^{3/2}}{(1 + 2\ell nb/a)(1 - f)} \quad (6)$$

where b is the drift tube radius, a is the beam radius, γ_o is the relativistic mass ratio for the electrons at injection, and $f = n_i/n_e$ represents any charge neutralization provided by positive ions. As indicated in Fig. 4a, a virtual cathode forms at the injection point with a depth approximately equal to the anode-cathode potential difference. The axial position of the virtual cathode downstream of the anode plane is usually on the order of the anode cathode gap, so that megavolt potentials at the virtual cathode are formed only millimeters away from the grounded anode. It is this very high electric field on the upstream side of the virtual cathode that collective accelerators usually seek to exploit for ion trapping and acceleration.

If the beam is injected into a localized gas cloud (Fig. 4b), ionization processes can quickly build up sufficient ion density to neutralize the electron beam space charge and the virtual cathode can move downstream to the edge of the gas cloud. This motion can result in the acceleration of a few ions to energies considerably higher than the depth of the potential well at the virtual cathode.

In order to control the motion of the virtual cathode over distances greater than a few centimeters, however, a means of providing an ionization channel whose axial extent can be controlled as a function of time is required (Fig. 4c). In the present experiments, shown schematically in Fig. 5, the ionization channel is generated by time-sequenced laser-target interactions. A laser pulse is separated into many approximately equal energy beams which are then optically delayed over different path lengths. The laser light then vaporizes and ionizes a target material on the drift tube wall, and ions drawn into the beam by the electron space charge at the beamfront provide the required time-sequenced channel of ionization to control beamfront motion.

First Generation Laser-Controlled Beamfront Accelerator Experiments. During the past two years, we have completed our studies of the first generation Laser-Controlled Beamfront Accelerator experiment. In this experiments the injected electron beam pulse was 900 keV, 20 kA, 30 ns, and controlled beamfront motion and accompanying ion acceleration were attempted over a distance of 50 cm. Results of these experiments were published in the Journal of Applied Physics **66** (7), 1 October 1989, p. 2894-2898, and the Proceedings of the 1989 Particle Accelerator Conference, p. 624. The major conclusion of these studies are as follows:

Figure 4: Conceptual View of the Laser Controlled Beamfront Accelerator.

Figure 5: Schematic of the Laser Controlled Beamfront Accelerator.

- i) Effective control of a relativistic electron beamfront by a laser produced time-sequenced ionization channel has clearly been demonstrated for two different accelerating gradients (40 MV/m and 90 MV/m). In addition, beamfront motion without the laser-produced ionization channel is consistent with theoretical expectations, as is the rapid propagation observed when the ionization channel is produced well in advance of electron beam injection.
- ii) Controlled collective acceleration of ions at a rate of 40 MV/m over a distance of 50 cm has also been demonstrated, but acceleration at the higher gradient (90 MV/m) over this distance has not been demonstrated to date. As discussed in the theoretical progress section of this report, both the successful acceleration of protons at the 40 MV/m gradient and the failure to accelerate ions at the higher gradient over the entire 50 cm distance are entirely consistent with numerical simulations of the beamfront accelerator and simple analytic theory. In the higher gradient experiments, it was simply not possible to maintain the required > 90 MV/m electric fields at the virtual cathode over the entire 50 cm accelerating distance.
- iii) Approximately 10^9 protons/pulse were accelerated to a peak energy of 18 MeV in the 40 MV/m gradient experiments. Data from stacked foil measurements indicate that the spectrum has a strong peak at about 18 MeV.

Second Generation Laser-Controlled Beamfront Accelerator Experiments. A second generation Laser-Controlled Beamfront Accelerator has been designed with the aid of numerical simulations described in the theoretical section of this report. The specific design parameters and objectives for the new experiment are discussed in the theoretical section. The second generation experiment is designed to accelerate protons over a 100 cm distance at gradients up to 60 MeV/m using electron beams with energies in the range 1.2–1.5 MeV. Construction of the new experiment was completed over the past year (Fig. 6) and initial experiments were conducted.

Optical system specifications and tests are displayed in Tables 1–4. Table 1 specifies the different laser channel path lengths and D_n (the distance between the fully reflecting and partially reflecting mirrors for the n th channel) for a one meter, two-part accelerating gradient designed to accelerate protons to a peak energy of 35 MeV. The two-part gradient was designed on the basis of numerical simulations to be discussed in section II.B. Tables 2 and 3 detail the theoretical and actual partial reflecting mirror specifications, predicted laser energy at each spot, and actual experimental values achieved. Reductions in actual laser energy achieved over theoretical expectations are due to optical losses in the system.

Figure 6: Photograph of new experimental laser controlled beamfront accelerator system.

Given the complex nature of an optical system employing 40 mirrors, the optical losses are no higher than should be expected. Table 4 specifies the mirror positions for a somewhat slower gradient designed to achieve a peak proton energy of about 25 MeV.

In order to monitor beamfront propagation down the drift tube, four wall current probes were inserted 11.5, 36.5, 61.5, and 86.5 cm downstream of the injection point to measure beam current density at the drift tube wall. Typical oscilloscope waveforms from beamfront propagation measurements are shown in Fig. 7. These waveforms exhibit characteristics consistent with a picture of the beam propagating in a well focused manner to the end of the ionization channel, and then rapidly exploding to the drift tube wall at the position of the virtual cathode (see Fig. 4c). In fact, it is encouraging to see that the wall probe current drops off quite rapidly as the beamfront passes by the probe position.

Plots of beamfront propagation distance vs time in the drift tube obtained from these waveforms are shown in Fig. 8 for various delay times between the firing of the laser and electron beam injection for the two-part gradient. It is easily seen that the designed beamfront control can be achieved if, as expected, the laser is fired immediately before the beam pulse. If the laser is fired too early, a plasma channel simply fills the drift tube in advance of beam injection and beam propagation occurs quite rapidly down the tube. Measurements of wall current peak magnitudes as a function of axial position are shown in Fig. 9 for variety of laser-beam firing delays. These results show, as expected, that the laser channel does allow for significant current propagation downstream compared to the no-laser results. Current levels observed far downstream, however, are consistently low. This is probably attributable to the fact that the pulse line accelerator used in the experiments has a usable pulse duration of only about 20 ns. Thus the injected electron beam pulse is probably already decreasing rapidly in magnitude by the time the beamfront is 85 cm downstream.

Figure 10 displays beamfront propagation data for the slower gradient mirror configuration discussed in Table 4. As can be readily seen, the designed beamfront propagation control has been achieved in this case as well. Wall probe current levels for this configuration are shown for various laser-beam firing delays in Fig. 11.

2. Theoretical Studies

We will discuss the progress made during the past three years in four separate but related topics that were under investigation. The four topics are: "Initial Ion Production and Acceleration Phase," "Laser Controlled Beamfront Accelerator," "Intense Beam Propagation Across a Magnetic Field," and "Ion Hose Instability of the Bennett Profile Beam." The last two topics form part of the Ph.D. Thesis by Dr. Xiaohao Zhang (support provided by this contract; present address, Brookhaven National Laboratory).

Table 1: Path lengths and D_n for the twenty channels in the laser system.

Channel	Path Length (m)	D_n (m)
1	0.587	0.587
2	0.975	0.427
3	0.644	0.319
4	0.485	0.255
5	0.391	0.213
6	0.330	0.183
7	0.287	0.161
8	0.259	0.141
9	0.241	0.126
10	0.225	0.115
11	0.214	0.107
12	0.206	0.101
13	0.197	0.096
14	0.191	0.092
15	0.185	0.087
16	0.182	0.085
17	0.177	0.082
18	0.173	0.079
19	0.169	0.077
20	0.166	0.075

Table 2: Theoretical mirror reflectivities and the calculated energy per channel: actual mirror reflectivities and the calculated energy per channel.

Channel	Theoretical Mirror Reflectivity	Calculated Laser Energy (Joules)	Actual Mirror Reflectivity	Calculated Laser Energy (Joules)
1	0.950	0.3	0.95	0.300
2	0.947	0.3	0.95	0.285
3	0.944	0.3	0.95	0.270
4	0.941	0.3	0.95	0.257
5	0.937	0.3	0.95	0.244
6	0.933	0.3	0.925	0.348
7	0.929	0.3	0.925	0.322
8	0.923	0.3	0.925	0.298
9	0.917	0.3	0.925	0.276
10	0.909	0.3	0.90	0.339
11	0.900	0.3	0.90	0.306
12	0.890	0.3	0.89	0.303
13	0.875	0.3	0.88	0.294
14	0.857	0.3	0.86	0.301
15	0.833	0.3	0.83	0.315
16	0.800	0.3	0.80	0.307
17	0.750	0.3	0.75	0.307
18	0.667	0.3	0.67	0.304
19	0.500	0.3	0.50	0.309
20	—	0.3	—	0.309

Table 3: Summary of optical system tests.

Channel	Axial Position (<i>cm</i>)	Partial Mirror Reflectivity	Theoretical Laser Energy (<i>Joules</i>)	Experimental Laser Energy (<i>Joules</i>)
1	5	0.95	0.300	0.30
2	10	0.95	0.285	0.28
3	15	0.95	0.270	0.265
4	20	0.95	0.257	0.25
5	25	0.95	0.244	0.23
6	30	0.925	0.348	0.315
7	35	0.925	0.322	0.295
8	40	0.925	0.298	0.275
9	45	0.925	0.276	0.255
10	50	0.90	0.339	0.30
11	55	0.90	0.306	0.25
12	60	0.89	0.303	0.24
13	65	0.88	0.294	0.24
14	70	0.86	0.301	0.23
15	75	0.83	0.315	0.225
16	80	0.80	0.307	0.22
17	85	0.75	0.307	0.21
18	90	0.67	0.304	0.195
19	95	0.50	0.309	0.20
20	100	—	0.309	0.19

Table 4: Path lengths and D_n for the twenty channels in the laser system.

Channel	Path Length (m)	D_n (m)
1	0.460	0.460
2	0.781	0.389
3	0.678	0.337
4	0.598	0.297
5	0.535	0.265
6	0.483	0.239
7	0.442	0.218
8	0.406	0.200
9	0.377	0.185
10	0.349	0.171
11	0.328	0.160
12	0.306	0.149
13	0.291	0.141
14	0.275	0.133
15	0.260	0.125
16	0.248	0.119
17	0.237	0.113
18	0.225	0.107
19	0.216	0.102
20	0.206	0.097

Figure 7: Typical traces of the charge/current probe used to measure the beamfront propagation.

Figure 8: Wall probe measurements of beamfront propagation versus time for different laser firing delay times and without the laser ("two-part" gradient.)

Figure 9: Voltage observed at the wall probes for different delay times ("two-part" gradient.)

Figure 10: Wall probe measurements of beamfront propagation versus time for different laser firing delay times and without the laser ("slower" gradient).

Figure 11: Voltage observed at the wall probes for the different delay times ("slower" gradient).

Initial Ion Production and Acceleration Phase During the second year, we completed numerical studies related to the initial phase of the ion production and acceleration process in a localized gas cloud system.

In experiments in which an intense relativistic electron beam is injected into an evacuated drift tube with a localized gas cloud located near the anode that serves as the source of ions for effective beam propagation, ions with energies several times the beam energy as well as a large fraction of the injected electron current are observed downstream. These experiments have been simulated using a particle-in-cell code which realistically models ionization of the gas. We found that when the injected electron beam current exceeds the space-charge limiting current, the majority of the ions produced achieve energies of the order of the beam energy and provide for an effective channel to space charge neutralize the electron beam. There are a few ions that are accelerated to energies several times the electron beam energy by the coherent motion of the ions and the intense virtual cathode electric fields. The majority of the ions allow for the total beam current to propagate to the downstream surface once the ion channel has also propagated to this location. The dependence of the peak ion energy on the system parameters as observed in the simulations was also examined. For the parameter regimes investigated, beam energies up to 3 MV, beam currents up to 35 kA, gas pressures up to 600 mTorr, and gas cloud widths up to 6 cm, peak ion energies of 5–6 times the electron beam energy have been observed in the numerous simulations.

Laser Controlled Beamfront Accelerator–Simulation. During the first and second years, we have written an electrostatic particle-in-cell code called LCA to simulate the laser-controlled acceleration experiments. A schematic of the simulation model is shown in Fig. 12. In the simulations, an electron beam of voltage V_0 , current I_0 , and radius R_b is injected into a grounded cylindrical drift tube of length d and radius R_w along the axis of the drift tube. The region of the drift tube extending from the anode to a distance z_0 downstream is filled by hydrogen gas at a constant pressure p_0 . The electron beam is assumed to be focussed by an infinitely strong guide magnetic field, so that particles in the simulation move only along the axis of the drift tube. The beam radius is also assumed to be much smaller than the wall radius so that the charge and current density and the axial electric field are approximately uniform across the beam cross-section.

The macroparticles in the simulation obey the relativistic equations of motion, the electric field E_z on axis is computed by first computing the potential ϕ at equally spaced grid points on the axis, numerically computing the derivative at points lying halfway between successive grid points, and linearly interpolating to obtain the value of E_z at other points. Ionization of the neutral gas is modeled by dividing the gas region into grid cells and monitoring the

Figure 12: Simulation model

amount of ionization in each grid cell which is produced by impact ionization. When the total number of ions produced in a grid cell exceeds the number of ions in an ion macroparticle, an ion macroparticle and an electron macroparticle are created at the center of the grid. The actual details have been explained elsewhere.

The time required for the laser beam to create plasma after striking the target is assumed to be a constant quantity which is specified. The front of the laser-produced plasma is assumed to sweep smoothly from one end of the drift tube to the other and the only effect of the plasma in the simulation is to completely neutralize any space-charge in the drift tube behind the plasma front.

The first set of results was obtained for a 900 keV, 20 kA, 1 cm-radius electron beam which is injected into a 5 cm-radius, 50 cm-long drift tube with a 2 cm-wide, 100 mTorr cloud of hydrogen gas located next to the injection plane. These parameters are those associated with the successful first generation laser-controlled beamfront accelerator experiments. As in the experiment, the front of the laser-produced plasma is assumed to travel down the drift tube at a velocity which increases linearly from $0.04c$ to $0.2c$ over a distance of 45 cm. Figure 13 shows the peak proton energy measured 45 cm downstream from the injection plane as a function of the time delay between the start of the beam pulse and the start of the laser pulse. For these runs it was assumed that the laser-produced plasma was created 10 ns after the laser beam struck the target and that the time required for the plasma to travel from

Figure 13: Peak proton energy measured at 45 cm versus time delay between start of laser pulse and start of beam pulse for linear velocity gradient.

the wall to the center of the drift tube was equal to the time required for a proton to travel from the wall to the surface of the electron beam, assuming that the potential depression produced by the beam was equal to the beam energy V_0 . The figure shows that the design energy of 18.8 MeV is attained over a broad range of time delays. If one compares the phase-space trajectory of the peak-energy proton macroparticle and the phase-space trajectory of the laser-produced plasma front, we find for this design that the accelerated proton tracks the laser beam trajectory closely.

When the length of the drift tube is double from 50 cm to 100 cm (with all other system parameters the same) and the same velocity gradient is used, the peak proton energy measured at the end of the drift tube falls short of the design value, i.e., the original velocity gradient cannot be extended to longer distances. The reason the original velocity gradient cannot be scaled to longer distances can be seen as follows. The equation of motion of a proton which is being accelerated by an electric field E_z is $dv/dt = (e/m)E_z$, which can be rewritten as $E_z = (mc^2/e)v dv/dz$. If the velocity gradient dv/dz used in the above runs is substituted into the preceding expression, we find that the electric field needed to accelerate the proton at the desired velocity gradient is approximately $E_z = 334 \beta$ MV/m, where $\beta = v/c$, e.g., for $\beta = 0.3$, an electric field of more than 100 MV/m is needed. In Fig. 14, we have plotted the magnitude versus the location of the peak electric field for the run with a 100 cm-long drift tube in which the greatest peak proton energy was measured. Notice that as the beam front moves downstream, the peak electric field tends to fall until it is no longer large enough to continue accelerating the proton at a constant velocity gradient.

The tendency of the peak electric field to fall as the beam front moves downstream suggests that it may be better to accelerate the protons with a steep velocity gradient at the start and then taper the gradient as the beam front moves downstream. In the second series of simulations, we therefore chose to do a set of runs in which the beam front velocity increases linearly with distance (i.e., $dv/dz = \text{constant}$) until it reaches a transition point, after which the beam front velocity increases linearly with time (i.e., $dv/dt = \text{constant}$).

The plasma front velocity at the anode plane was chosen to be $0.04c$, the plasma front velocity at the downstream end of the drift tube was chosen to be $0.4c$, and the velocity $\beta_t c$ at the transition point z_t was varied. Figure 15 shows the peak proton energy measured at 90 cm as a function of β_t for three different values of z_t . Note that $V_0 = 1.5$ MV in this run. The figure shows that by adjusting β_t , energies in excess of the design energy of 60.7 MeV can be achieved.

It should be noted that, as might be expected, the peak proton energy depends strongly on the beam energy, since the peak electric field increases with increasing beam energy. In Fig. 16, we plot the peak proton energy measured at the downstream end of a 50 cm-long

Figure 14: Magnitude versus location of peak electric field.

Figure 15: Peak proton energy measured at 90 cm versus transition velocity β .

Figure 16: Peak proton energy measured at 50 cm versus electron beam energy.

drift tube as a function of the electron beam energy. In each case the same velocity profile is used. For beam energies greater than 1.4 MeV, peak proton energies exceed the design value. As the beam energy decreases, the peak proton energy decreases sharply, e.g., at 1 MeV, the peak proton energy is already less than 20 MeV.

Optimal Acceleration—New Design. The largest proton energies achievable in any given system can be calculated by integrating the peak electric fields obtained from the simulations as a function of distance (assuming that the magnitude of the peak electric field depends only on the position of the plasma beam front and not on its velocity). The required variation of the plasma front position with time can be obtained by integrating the equations $v dv/dz = eE_{z,\text{max}}(z)$ and $dz/dt = v(z)$, where $E_{z,\text{max}}(z)$ is the peak electric field at z . Figure 17 shows β vs. z and t vs. z for a 1.2 MeV, 20 kA electron beam which is injected into a 50 cm-long drift tube.

We have developed the electrostatic PIC code LCA to simulate beam propagation and collective ion acceleration in a laser-controlled accelerator. Although the code only crudely models the experiment, it has been able to reproduce some of the experimental results and may be useful as a design tool for future laser-controlled acceleration experiments.

Figure 17: Optimal acceleration profile.

Ion Hose Instability of the Bennett Profile Beam. During the second and third years, we analyzed the ion hose instability of the self-consistent Vlasov equilibrium that leads to the Bennett radial density profile for the downstream electron-ion beam system. Specifically, we consider the limiting case where the ion beam species is cold, stationary, and provides charge neutralization. The most simple model assumes that each beam is rigid. We then investigate the effects due to the anharmonic nature of the potential well using a “spread density” model and the effects due to an axial thermal velocity spread of the electron beam. In addition, nonlinear effects due to large beam displacements are numerically investigated.

For the rigid beam model, the linearized equations for transverse motion of the beams are

$$\ddot{x}_e = \frac{\omega_e^2}{6}(x_i - x_e)$$

$$\ddot{x}_i = \frac{\omega_i^2}{6}(x_i - x_e)$$

and $\dot{z}_e = v_0$, $\dot{z}_i = 0$. We have defined $\omega_e = \sqrt{e^2 n_{i0} / \epsilon_0 \gamma m_e}$, $\omega_i = \sqrt{e^2 n_{e0} / \epsilon_0 m_i}$, $n_{i0} = n_{e0} = n_0$, and the displacements x are assumed to be much less than the beam cross-sectional dimensions a . For harmonic displacements, we find that resonant growth occurs for wavelengths

$$kv_0 = \frac{2\pi}{\lambda} v_0 = \frac{\omega_e}{\sqrt{6}} = \frac{2\pi}{\lambda_{\beta_e}} v_0$$

with a complex frequency

$$\omega = \left(\frac{1}{2} + i \frac{\sqrt{3}}{2} \right) \left(\frac{\omega_i^2 \omega_e}{12\sqrt{6}} \right)^{1/3},$$

where the growth rate is $Im\omega$. The spectrum of growth rates is shown in Fig. 18 by the solid line (R.B.), where the strong resonant growth is clearly seen. The system parameters for this case are $V_0 = 1$ MV, $I_0 = 20$ kA, $\beta_0 = v_0/c = 0.65$, $a = 1$ cm, and $R_w/a = 10$.

Figure 18: Ion hose instability growth rates versus wavenumber. System parameters are: $V_0 = 1 \text{ MV}$, $I_0 = 20 \text{ kA}$, R.B. - Rigid beam model, S.D. - Spread density model (anharmonic potential well), T.V. - Thermal velocity spread model.

The effects due to the Bennett profile which results in particles oscillating at different transverse frequencies are modelled by annular layers with different densities. After weighting by density and averaging to obtain the net beam displacement, we obtain a dispersion relation

$$1 - \left[1 + \frac{\omega}{\sqrt{2}\omega_i} \ln \frac{\omega - \omega_i/\sqrt{2}}{\omega + \omega_i/\sqrt{2}} \right] \left[1 + \frac{\omega - kv_0}{\sqrt{2}\omega_e} \ln \frac{\omega - kv_0 - \omega_e/\sqrt{2}}{\omega - kv_0 + \omega_e/\sqrt{2}} \right] = 0.$$

These results are displayed in Fig. 18 by the dashed curve labelled S.D. We see that the spread in transverse oscillation frequency (weighted by density) results in an unstable spectrum up to the peak on axis-value of $\omega_e/\sqrt{2}$, versus up to the average value of $\omega_e/\sqrt{6}$ for the rigid beam (R.B.) model.

The effects due to an axial thermal velocity spread of the electron beam species is modelled by the one-dimensional relativistic Maxwellian distribution. The beam is assumed to be composed of many rigid disks having the same density profile but different axial velocities. These results are shown by the dashed curve in Fig. 18 labeled T.V.

In summary, our studies of the ion hose instability of a Bennett profile beam indicate:

- From the rigid beam model, the instability exists in the long wavelength region and is weakly absolute. The resonant growth occurs at $\lambda = \lambda_{\beta_e}$.
- Due to the anharmonic potential as well as axial velocity spread, the maximum growth is reduced, the unstable region is broadened, and the instability becomes convective.
- For the case of large beam displacements, the nonlinear equations of motion for the ion hose oscillations have been derived for the Bennett profile. As the beam separation becomes of the order of the beam radius, the instability growth decreases and saturates.

[This work was part of X. Zhang's Ph.D. thesis.]

Intense Beam Propagation Across a Magnetic Field. During the second year, we completed the studies on a model of the propagation of an intense electron-ion beam across a transverse magnetic field.

Previous theoretical studies have shown the existence of a self-consistent downstream Bennett equilibrium for the electrons and ions when no applied magnetic field is present [C.D. Striffler, R.L. Yao, X. Zhang, Proc. of the 1987 IEEE PAC, page 975]. We have related these downstream properties to the diode voltage, the transmitted electron beam current, and the ion properties in the localized gas cloud region. For our experimental parameters, $V_0 = 1$ MV, $I = 5$ kA, the downstream self-pinched equilibrium state is composed of cold ions with low axial speed (< 0.05 c), and electrons with a temperature of about 60 keV and an axial

speed of about 0.84 c. The nearly charge-neutral beam system has been shown to effectively propagate up to the diode current of 20 kA where the electron temperature is predicted to be about 200 keV. [This self-consistent equilibrium model was completed during the first year and was part of X. Zhang's Ph.D. thesis.] The 5 kA case was chosen for examining the effects of a transverse magnetic field on the electron-ion beam system, mainly because of experimental reproducibility. We considered the propagation of this intense electron-ion beam across an applied magnetic field. We find that in the intense beam regime, the propagation is limited due to space-charge depression caused by the deflection of the electron beam by the transverse field. This critical field is of the order of the peak self-magnetic field of the electron beam which is substantially higher than the single particle cutoff field. [This work was completed during the second year and was part of X. Zhang's Ph.D. thesis.]

B. Study of Pseudospark-Produced Electron Beams

We have continued and expanded our experimental program on pseudospark discharges. Although pseudospark discharges have interesting switching characteristics which are being studied at other laboratories, our current interest is centered around the study of high brightness electron beams produced in such discharges. A simple discharge chamber of modular type was constructed to investigate this new electron beam source. The chamber was operated at a low voltage ~ 25 kV, producing ~ 10 Hz, electron-beam pulses of ~ 100 A, ~ 10 ns. The rms emittance of the resultant electron beam was measured and found to be ~ 65 mm-mrad, which corresponds to a normalized brightness of the beam of $\sim 2 \times 10^{10}$ A/(m²rad²). Results of these experimental studies were published in *Applied Physics Letters*, **56**, 1746 (1990) and *Physics of Fluids B* **2**, 2487 (1990). Recently, preliminary empirical scaling studies of breakdown voltage, beam current, and beam emittance have been carried out on a new experimental setup with an operating voltage of up to 50 kV. The breakdown voltage curve, shown in Fig. 19, is approximately a function of product $p\sqrt{d}$ unlike the conventional Paschen curve which is a function of Pd . The electron current increases with breakdown voltage up to ~ 25 kV and then tends to decrease, as shown in Fig. 20. The beam current and its duration increase with the external capacitance.

C. Emittance Measurement by a Finite-Width Slit

In actual emittance measurements, use of finite-width slits is inevitable for the following reasons: a narrower slit reduces the beam intensity resulting in a poor signal to noise ratio of the data; the minimum slit size is limited by the fabrication technique employed; a narrower width in a finite thickness slit plate reduces the maximum acceptance angle. The finite-width

Figure 19: Measured characteristic breakdown curves for the pseudospark chamber for various intermediate gaps. Here N corresponds to the number of insulator disks/gaps.

Figure 20: Variation of the peak electron beam current with the breakdown voltage for various values of intermediate gaps.

slit, in contrast to an ideal thin slit, introduces an error in phase-space measurements which in turn, easily causes a substantial error in emittance analysis. We have derived a numerical method of correcting for the effects of a finite-width slit to improve emittance and brightness evaluations. This study was detailed in a paper, "Phase-Space Measurement of a Beam with a Maxwellian Transverse Velocity Distribution by a Finite-Width Slit," submitted to Review of Scientific Instruments (October 1990).

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I. Introduction and Synopsis

Prior to the present work, stable gyrokystron amplifier operation was exemplified by the state-of-the-art development project at the Naval Research Laboratory (NRL) where output power of 50 kW in 1 μ sec pulses was realized at a frequency of 4.5 GHz. During this three-year research contract period, we aimed at examining the very ambitious proposition of whether gyrokystron power might be enhanced almost one thousand times (up to 36 Megawatts) while at the same time doubling the frequency to X-band and maintaining efficiency in the 30-45% range. This could not be accomplished by a simple scaling since the NRL gyrokystron operated in the fundamental mode in the cavities with all modes cut off in the drift spaces; such a conservative design would not handle the much higher power and higher frequency of interest in the present study. Thus, difficult problems involved in stabilizing overmoded structures had to be addressed.

Stable operation was achieved in an overmoded gyrokystron circuit with output power in the 2 to 3 megawatt range at 9.85 GHz; pulse length was approximately 1 μ sec. The efficiency was 6% but this represents the linear or unsaturated value; the gain was limited because only a two cavity gyrokystron circuit was tested initially. We built and cold tested a three cavity circuit to allow us to drive the amplifier into saturation. We also continued to improve our instability suppression techniques to allow extension of the operating parameter from a voltage of 425 kV to the design value of 500 kV, and from a velocity ratio of $\alpha = v_{\perp}/v_z = 1.1$ to the design value of 1.5.

Aside from increasing gyrokystron power by a factor of 50, a number of scientific contributions of general value were made in the course of this gyrokystron research; these contributions include the following:

1. techniques for loading circuits with lossy dielectrics to achieve stability were extended beyond the state-of-the-art;
2. diamagnetic loops were used for the first time to confirm the measurements of α obtained with electrostatic probes;
3. the phase noise variation with voltage in gyrokystrons was shown both theoretically

and experimentally to be of a dramatically different form from the variation in conventional klystrons;

4. the phase noise level in a 500 kV gyrokystron was shown to be compatible with requirements of the linac collider application especially if some feedback stabilization is employed;
5. the CASCADE code was developed to compute fields in cavities of complicated geometry which are loaded with lossy dielectrics (this is a unique computational capability); and
6. the QPB code was developed using the Lie transform method to compute start oscillation currents for the fields produced by CASCADE. (QPB is 50 times faster than previously available codes.)

II. Experimental Research Progress

During this contract period, we completed our studies of the Magnetron Injection Gun (MIG) performance and initiated stability and amplification studies in several two-cavity gyrokystron configurations. The experimental facility functioned well, with the modulator firing approximately a million times (at 1 Hz) for the gyrokystron studies alone. The beam studies verified simulated MIG performance over a wide range of parameters. In the gyrokystron tubes, significant improvements were made in suppressing the typical gyrotron instabilities that have occurred. To date, over 3 MW of amplified power with 18 dB gain and an efficiency of 6.5% (unsaturated) were achieved in 1 μ sec long pulses at reduced voltage and current levels. This power level represents an increase of a factor of 50 in the state-of-the-art of gyrokystrons with thermionic cathodes.

A. Facility Improvements

We installed remote controls both for the magnet supplies and the modulator voltage level. We added the capability to quickly download oscilloscope traces to the computer. We also designed, constructed, and installed an inductor on the bottom leg of the resistive divider (see Fig. 1) to minimize deviations in the control anode to cathode voltage ratio (V_{ca}/V_c). The need for this was discovered during the beam characterization tests, and is discussed in the next section. An RLC network was designed and installed across the thyratrons to eliminate the punchbacks that occurred near 500 kV. Our vacuum system was finalized with a total pumping capability of 240 ℓ /s in four units. This allowed us to safely operate up to the maximum machine repetition rate (5 Hz). The facility was certified radiation safe up to 400 kV at full current.

In our vacuum tube processing facility, we fabricated the vacuum jackets, microwave windows, and lossy ceramics required for our tubes. This in-house capability has significantly reduced the down time between configurations.

We increased significantly our microwave diagnostic capability (see Fig. 2). The output waveguide frequency measurement range was expanded to include J-band. Wave-guide runs were installed so that we could get input and output frequency information in the screen room

Figure 1: Voltage divider circuit for MIG.

on an HP8566B spectrum analyzer. A microwave horn was placed behind the electron gun to locate instabilities in the beam tunnel before the tube. Our water load and our directional coupler were calibrated. The bench test results showed that the coupler's response was flat from 9.75 to 9.95 GHz to within ± 0.5 dB. Reverse power was immeasurable with our network analyzer. Finally, we constructed a larger anechoic chamber with remote control positioning (translation and 90° rotation) of the pick-up antenna to facilitate mode pattern analysis.

B. Electron Beam Studies

Our investigation of electron beam characteristics focused mainly on the average MIG properties. We increased space-charge limited operation to 375 kV and 216 A. Additional parameter space searches continued to confirm the predictions of the e-gun code simulations. In addition to the studies on the second beam diagnostics configuration [see 1989 Progress Report], we installed and completed studies on a final configuration which had a new linear uptaper constructed with alternating rings of metal and carbon-impregnated alumino-silicate. Stability of the third configuration to microwave oscillations proved only marginally better than the previous configuration. This result supported our claim that the unstable modes exist

Figure 2: Microwave diagnostic hardware.

near cutoff in the beam diagnostics chamber (Fig. 3) so that the beginnings of the linear uptapers act as strong reflectors. Further improvements in stability could have been made, but were abandoned in favor of amplification studies.

Subsequent data analysis resulted in several interesting conclusions. First, sectorized capacitive probes are a viable means of estimating average velocity ratio and beam off-centering. Also, internal air-core current transformers are adequate alternatives to electrical breaks in this parameter range. Furthermore, diamagnetic loops yield velocity ratio information consistent with capacitive probe data. Finally, time-resolved capacitive probe data indicated that the velocity ratio would change by as much as 10% during the pulse "flat top". The root cause was traced to an improperly compensated resistive divider which resulted in a variation of V_{ca}/V_c . Our modulator performance codes were modified to allow an inductor to be simulated on the lower leg of the resistive divider. Optimal compensation was found to depend on beam voltage and current, but an inductance of $450 \mu\text{H}$ was a good nominal value. This inductor was constructed and tested during the amplifier studies and was shown indirectly (through microwave output power) to reduce velocity ratio variations considerably.

Figure 3: Beam diagnostics chamber.

C. Two-Cavity Experiments

The focus of the experimental studies was on the stability and amplification properties of two-cavity gyrokystron tubes. The nominal resonant frequency was 9.85 GHz and the quality factors were in the range 150-175. Four configurations were completely tested. The instabilities in the first tube were so severe that very little information was discerned. As the tubes progressed, we systematically increased the loading in various areas of the tube to improve stability. This culminated in tube IV, which had a stable parameter space covering beam voltages up to 425 kV, magnetic field variations of over 20%, magnetic compressions in excess of 11, and beam current exceeding 200 A.

Tube I. The core of tube I, which includes the input cavity, drift space, and output cavity is shown in Figure 4. The lossy rings in the drift tube were MgO with 1% SiC. The lossy rings before (and in) the input cavity were carbon-loaded porous alumino-silicate. The core's relation to the rest of the tube is depicted in the general schematic of Fig. 5. The output waveguide consisted of a straight 12 cm section, a non-linear taper, a 100 cm long beam dump, a second non-linear taper, and a half-wavelength output window. The beam

Figure 4: Two Cavity Gyroklystron (tube II).

tunnel downtaper before the input cavity was the same one used in the final two beam experiments and is shown in Fig. 6. The downtaper proved to have insufficient loss so that low beam power level instabilities tended to disrupt the beam in the compression region and cause pressure rises which would automatically shut down the system.

Tube II. The second tube retained the core of tube I but had a modified downtaper which replaced the internal current transformer space with more lossy rings. This configuration is referred to as the "previous design" in Fig. 7. The oscillation-free parameter space was still rather small, but enabled us to characterize the instabilities and try some amplification studies. There are four basic types of microwave instabilities. The first class of modes exist in the output waveguide in frequency ranges where the output window is a good reflector. These modes often occur in groups spaced by about 60 MHz, which corresponds to different axial mode numbers. Because the beam-microwave interaction occurs after the output cavity, these modes are usually eliminated by amplifier operation and are of little concern. The second class of modes exist in the straight section adjacent to the output cavity (with some of their energy apparently extending into the cavity). These modes are near cutoff and use the first nonlinear taper as a reflector. The most troublesome mode had a TE_{12} distri-

Figure 5: Gyroklystron Schematic

Figure 6: The new linear downtaper with lossy ceramic rings.

Figure 7: Linear downtaper designs

bution and a frequency in the range 9.8–9.9 GHz. These modes did interfere with amplifier operation, but were eliminated in tubes IV and V. The third type of mode is the “whole tube” mode. These modes are described in the theory section of this and previous progress reports. Basically, the energy is contained in the drift section and the cavities act as reflectors. These modes existed at several frequencies with a particularly disruptive one slightly above 9 GHz. These oscillations were of concern because they appeared to occur at current levels below theoretical predictions. The problem was traced to inadequate dielectrics and was corrected in tubes III–V. The final mode class involves oscillations in the beam tunnel downtaper. These modes are the most disruptive to tube operation and occur in two bands: 7.0–7.5 GHz and 8.1–8.4 GHz. These instabilities were identified as $m = 1$ modes.

A stable region for amplifier experiments was found at 80% of the nominal magnetic field at a beam voltage of 175 kV and a current of 55 A. Numerical simulations estimated the velocity ratio to be $\alpha \approx 0.45$ and the axial velocity spread to be about 4%. Our computer simulations of amplifier operation indicated that the gain would be about 0 dB under these restricted beam parameters. The experiment confirmed this prediction and output power was thus limited by the input source. Typical input and output crystal detector responses

Figure 8: Crystal detector response during amplifier operation. The upper trace shows the magnetron input power; the lower trace, the gyrokystron output power. Beam voltage is 175 kV, beam power, 9.6 MW, magnetic field, 0.452 T (80% nominal), magnetic compression, 9.9, velocity ratio, 0.45, and velocity spread, 4%.

are shown in Fig. 8. The lower trace is the output signal and its reduced width corresponds to the duration of the voltage pulse. A typical gain vs. frequency plot is shown in Fig. 9.

Tube III. The third tube featured a new drift tube (Fig. 10) and a new downtaper (labelled "current design" in Fig. 7). The new drift tube utilized a non-periodic configuration with tapered ceramic rings. The new lossy material was BeO with 20% SiC. Cold-test transmission data for the TE_{11} mode in X-band through the two drift tubes is shown in Fig. 11. While the original drift tube was quite lossy at the design frequency, there were two broad regions in X-band alone where the rings exhibited virtually no loss. Part of the trouble originated from incorrect (published) dielectric constant data. While the berrylia data was also wrong, the new drift tube still enjoys considerable broadband loss. TE_{11} transmission data from 7–9 GHz for the various downtapers is shown in Fig. 12. As indicated, significant progress was made in the attenuation through this critical component. The third and fourth downtapers are longer than the first two and cover a cavity formed by the body of a gate

Figure 9: Gain vs. frequency for 27 kW pin.

valve (which is used to maintain gun vacuum during the tube changes). The net result of these changes was a dramatic increase in the stable parameter ranges due to suppression of the downtaper and whole-tube instabilities. The resonant frequencies and Q 's of the two cavities were slightly affected by the new drift tube and were readjusted.

The enhanced stability allowed amplification studies up to 350 kV. Output powers in excess of 1.6 MW were achieved at 305 kV and approximately 108 A. The gain was about 15 dB and the efficiency was near 5%. The optimum magnetic field and input frequency were 0.452 T and 9.866 GHz, respectively. Higher magnetic fields could not be used because the second type of instability would persist for the full duration of the voltage pulse. The magnetic compression was limited to 10.5 by instabilities in the downtaper. The resulting perpendicular to parallel velocity ratio was estimated to be about $\alpha = 0.9$ from simulations, which is significantly below the design value of $\alpha = 1.5$. Again, the amplifier's performance was in reasonably good agreement with the theoretical predictions. Figure 13 shows gain as a function of frequency for various beam currents. Note that the lower currents peak near 9.85 GHz while higher currents tend to peak near 9.87 GHz. This pulling effect appears to be quite standard. Another interesting effect is shown in the output power vs. input drive

Figure 10: Two Cavity Gyroklystron (tube III)

Figure 11: Drift tube attenuation.

Figure 12: Measured attenuation in all downtaper configurations.

power of Fig. 14. Several of the constant current curves show steep increases in gain, which we believe is due to beam modification of the EM field profiles (an effect not included in our theoretical models but found consistent codes). Figure 15 shows the spectrum for the optimal point. Figure 15c shows the frequency spectrum of the TE_{12} mode in the output section when the drive signal was off. This mode persisted when the drive signal was off, but, due to slightly varying voltage conditions, was restricted to the beginning of the pulse. Figure 15a shows the spectrum when the magnetron was on. There is clearly maximum energy at 9.866 GHz due to the amplification which occurred late in the pulse. There were also downtaper modes present in the range 6.8–7.5 GHz, as shown in Fig. 15b.

Tube IV. The only difference between tubes III and IV is in the output waveguide adjacent to the second cavity. Since the second class instabilities were near cutoff, a 2° taper was introduced to the waveguide wall. This made the radius 0.4 cm smaller just after the output cavity so the coupling aperture was adjusted to maintain the proper cavity Q . This modification totally eliminated the second class of instabilities and allowed us for the first time to study amplification at magnetic fields approaching the design values. The limitation due to downtaper modes remained, and we were still limited to reduced values of α . Figure

Figure 13: Tube III gain relative to drive power at various beam currents.

Figure 14: Tube III power output vs. drive power at various beam currents.

Figure 15: Selected frequency spectrum for typical Tube III operation.

Figure 16: Tube IV power output vs. frequency with 53 kW drive power.

16 shows the nominal output power vs. frequency results for a range of magnetic fields. As indicated, powers in excess of 2 MW are easily obtained at 95% of the full field value of 0.565 T. Figure 17 shows the output power vs. drive power for the various magnetic field settings. Again the gain enhancement phenomena exists. Figure 18 shows the dependence of gain on magnetic compression (and hence α). Note that none of the curves are saturated. This implies that if the downtaper modes could be further stabilized, higher powers could be obtained even at these reduced beam power levels. The nominal input cavity coupling efficiency is shown in Fig. 19. The nominal coupling loss is ≈ 3 dB. As expected, the beam modifies coupling with the best enhancement occurring at maximum gain.

At 85% of full field, there was a narrow parameter range that gave output powers in excess of 3 MW. The pulse length for this enhanced operation was only a few hundred nanoseconds, but only because the ideal conditions (higher α) didn't persist for the entire voltage pulse. The comparison of the output power vs. frequency for the enhanced and normal modes are shown in Fig. 20.

The burst in α due to improper compensation of the resistive divider was mentioned before in the beam section. While it helped to produce the 3.1 MW power level, it nominally

Figure 17: Tube IV power output vs. drive power at 9.87 GHz.

Figure 18: Tube IV gain relative to drive power vs. compression at 9.87 GHz.

Figure 19: Tube IV coupling vs. frequency.

Figure 20: Tube IV power output at 85% B_{nom} and 51.6 kW drive power.

hurts operation by enhancing growth of the downtaper modes. In Fig. 21 we display the beam voltage pulse (top curve in *a*) and five microwave detector signals. The curve in Fig. 21a labelled "X-band" can sense signals between 6.56 GHz and 14 GHz. The choppy signal on the first half of the pulse represents the downtaper modes. This corresponds to the lower trace in Fig. 21c, which shows the microwave signal at the back of the gun. The second half of the pulse represents the amplified signal. It would be considerably wider if not for the presence of the other mode. Fig. 21b shows the relative incident and reflected powers at the input cavity window. The second spike on the reverse signal is not reflected power but rather an indication that energy from the downtaper modes has entered the input cavity. The "Ku-detector" on the upper curve of Fig. 21c actually registers everything above 9.49 GHz. The two early pulses represent harmonics of the downtaper modes and the latter pulse represents the amplified signal.

Figure 22 shows a set of microwave signals for a similar case after the new inductor was installed and the voltage pulse was flattened. Note that the downtaper modes were completely eliminated (though they would return at higher compression). Now the flat top region is essentially 1 μ sec. Figure 23 shows typical spectra for the final configuration of tube IV. Comparison with Fig. 15 shows dramatically the improvements we have made in tube stability.

Considerable work was done to verify the power levels. One such effort was to identify mode purity via radiation patterns. Figure 24 shows the theoretical radiated pattern as the solid line. The experimental points (circles) are in excellent agreement. The zero readings on axis (squares) are for E_θ pickup and indicate pure TE_{01} 2 MW signal.

Tube V. The final two-cavity tube has two modifications. First, the output cavity Q was raised (by 225) in attempt to push the tube to saturation at the reduced power levels. Second, the downtaper region was modified to increase TE_{11} transmission loss. The cold-test TE_{11} transmission results are displayed in Figure 12. Notice that the critical range 7-7.5 GHz shows an additional 6 dB loss. Finally, the lossy rings just before the input cavity have also been adjusted to increase loss. Experimentation is in progress.

Figure 21: 90% 352 kV $f_M \sim 11$ B_{nom} 127 A 9.86 GHz

Figure 22: With 95% 407 kV 11.1 f_M inductor B_{nom} 134 A 9.87 GHz

Figure 23: Selected frequency spectrum for typical Tube IV operation.

Figure 24: Measured and computed power pickup vs. angle from axis.

D. Three-Cavity Design and Cold Testing

The experimental effort to date, as regards RF generation, was devoted to studying the two-cavity gyrokystron circuit. That effort has achieved the elimination of many unwanted modes and has achieved multimewatt operation. We are now ready to test a three-cavity system which is expected to achieve higher gain and somewhat higher efficiency.

Figure 25 shows the three-cavity system and Table 1 shows the full power operating parameters predicted by our nonlinear amplifier code. Except for RF coupling and tuning connections, the input and buncher cavities are identical, each has an outer radius of 3.1 cm and is 1.68 cm long. The outer wall of each cavity is lined with 2 mm of lossy dielectric to achieve the desired Q and to suppress unwanted oscillations. The drift tube on each side of the cavities has a short section of metal to isolate the cavities from the microwave absorbers. The output power is extracted axially from the output cavity. There is no dielectric loading in the output cavity and the Q is controlled by the height and thickness of a lip at the end of the cavity.

The microwave circuit is designed in two steps: (1) The geometry and Q of the cavities are optimized for gain and efficiency, and (2) the drift tubes are loaded with lossy dielectrics

Figure 25: Schematic diagram of three cavity circuit.

Table 1: Beam and operating parameters predicted by our nonlinear amplifier code.

Beam:	Power	80 MW
	Voltage	500 kV
	Velocity Ratio (α)	1.5
	Velocity Spread	6.8%
	Center Radius	0.79 cm
	Guiding Center Spread	0.268 cm
	Axial Magnetic Field	0.5–0.6 T
System:	Frequency	9.85 GHz
	Input Cavity Q	225
	Buncher Cavity Q	225
	Output Cavity Q	170
	Theoretical Efficiency	42%
	Theoretical Gain	60 dB

to prevent unwanted oscillations.

The use of lossy dielectrics to suppress unwanted modes appears simple. There are, however, a number of considerations. The two most important are choosing the real and imaginary parts of the dielectric constant and deciding on a configuration, e.g., alternating rings of dielectric and metal versus solid dielectric. For high power tubes, the knowledge of the dielectric properties over a rather large frequency band, in our case 6–20 GHz, is essential. Thus, the issues of measuring both the real and imaginary parts of the dielectric constants and modeling the system numerically are important.

We have recently completed a study of seven vacuum compatible, lossy, ceramic dielectrics [see Appendix]. We placed samples of each dielectric at the end of a shorted waveguide. We used the complex reflection coefficient to calculate the complex dielectric constant over the frequency range 8–12.4 GHz. In the near future, we will increase the frequency range of our measurements to 6–26 GHz with our new microwave source.

To model the dielectric loaded complex circuit, we have developed a code based on the scattering matrix method.[1] Our code, "CASCADE", assumes cylindrical symmetry and allows one radial region of lossy dielectric material (e.g., vacuum for $r < r_0$ and dielectric for $r_0 < r < r_w$ where r_w is the outer wall radius). In addition, we have a code that calculates the start oscillation current in a complex dielectric loaded circuit.

Preliminary beam tests on the two-cavity system revealed that unwanted oscillations occurred from 6 to 20 GHz. Nonresonant absorbing structures in the drift regions were shown to improve the suppression of modes for the two cavity system (see section II.C). To achieve attenuation over the range 6 to 20 GHz we are considering a slightly different nonresonant absorbing system for the three-cavity circuit.

The operating mode in all three cavities of the gyrokystron is the TE_{011} cavity mode. To isolate the cavities, the TE_{01} mode attenuation in the drift section between cavities must be greater than the intercavity gain, which is 40 dB between the input and buncher cavities and 20 dB between the buncher and output cavities. To achieve this attenuation in the nonresonant system we place a smooth metal drift section 1.5 cm long on each side of the input and buncher cavities and on the upstream side of the output cavity; see Fig. 25. The remaining space between the cavities contains the absorbing lossy dielectric. The metal

regions isolate the TE_{011} mode at 9.85 GHz in the cavities from the absorbing structures. At 9.85 GHz the TE_{01} mode is evanescent in these regions and is isolated by 12 dB/cm. Thus, in the first drift region the necessary isolation of 40 dB plus a 10 dB safety margin can be achieved if the 4.25 cm dielectric region has attenuation greater than 3.3 dB/cm.

For other than the operating mode, the absorbing structure is designed by optimizing the attenuation in the least attenuated mode. Depending on the frequency and kind of dielectric used, the least attenuated mode can be one of a number of hybrid modes which exist in a lossy dielectric lined tube. In practice we measure the attenuation in the least attenuated mode by injecting the TE_{11} mode.

To determine the best absorbing system to use in the three-cavity circuit drift regions we measure the power absorbed in the TE_{11}^o and TE_{01}^o modes by a number of different absorbers. We use an X-band network analyzer and pairs of $TE_{10}^o \rightarrow TE_{11}^o$ and $TE_{10}^o \rightarrow TE_{01}^o$ waveguide mode converters to generate the test signals. The two best absorbing systems we found are a drift tube lined with 3mm of carbonized aluminum silicate (CAS)[2] and a drift tube lined with a stack of rings made of beryllia mixed with 20% silicon carbide (see Section II.C). Each of the beryllia rings has an outer diameter of 6.2 cm and is 0.742 cm long. The inside radius of the rings are tapered at 45°. Figure 26 shows a diagram of the two absorber systems and Fig. 27 compares the TE_{11} mode attenuation of each system in X-band (8–12.4 GHz). Both systems give sufficient attenuation to isolate the cavities in the operating mode. Each system has some advantages. The beryllia system is less porous and will outgas less in the vacuum system, and the CAS system gives better attenuation in X-band. We will use the CAS system in our first three-cavity circuit. Figure 28 compares the measured attenuation to the attenuation calculated by CASCADE for the TE_{11} mode in a drift tube with a 3.75 cm CAS liner as in Fig. 26b.

The cavity design for the gyrokystron evolves from our nonlinear design code. This code allows us to optimize the gain and efficiency by adjusting the circuit geometry and magnetic field. The quality factors of the cavities are set relative to start oscillation current in the operating mode. The magnetron input source injects RF into the input cavity through a single axial slit in the radial wall.

To optimize the gain of the three-cavity system through penultimate tuning it is necessary

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Figure 26: Axial cross-section of two drift tube absorber systems. The tube in (a) is made of a stack of tapered BeO + 20% SiC rings and one metal ring, and is used in the present two cavity circuit. The tube in (b) is a simple liner of carbonized silicate (CAS) and will be used in the first three cavity circuit.

Figure 27: Comparison of drift section attenuation for 2 mm carbonized aluminum silicate liner (solid line) and for 1.5 cm BeO + 20% SiC liner (dotted line).

Figure 28: Comparison of scattering matrix code "CASCADE" with measurement, for the 3 mm thick carbonized aluminum silicate absorber.

to adjust the resonant frequency of the buncher cavity. The tuning is achieved via a pair of 1/8 inch OD metal rods which can be inserted into the buncher cavity (Fig. 25). The rods are positioned by a pair of micrometer linear motion feedthroughs located just outside the magnetic field coils. Figure 29 shows the relation between rod position and cavity resonant frequency. Except for these coupling slots the input and buncher cavities are identical.

Because the system operates at such high current, dielectric loading in the cavities is necessary to prevent spontaneous oscillation. The SOC code is used to design the cavity loading. The dielectric loading of the cavity must discriminate between the TE_{011} and TE_{021} modes. That is, the Q of the TE_{021} mode should be considerably reduced while leaving the Q of the TE_{011} mode high enough for efficient bunching of the beam. The structure which we found that works the best is a simple ring of lossy dielectric placed on the radial wall of the cavity. To operate at 70% of the start oscillation current with a beam of 160 A requires a cavity Q of 320. Figure 30 shows the computed start oscillation current of different modes after loading the cavities with dielectrics. As shown, the start oscillation current is well above 200 A for the operating mode. We are in the process of computing the start oscillation currents for the whole tube modes.

Figure 29: Buncher cavity tuning curve. Negative insertion indicates that the tip of the tuning rod is outside the cavity.

Figure 30: Results of SOC for the input and buncher cavity design. The TM₀₁₀ mode is not unstable in this region of magnetic field. The operating current is 160 A.

E. Phase Noise Reduction in a Gyroklystron Amplifier

An expression for phase fluctuations in the output signal of a gyroklystron due to power supply voltage ripple was obtained in a form that clarifies the parametric dependence of the fluctuations, i.e.,

$$d\phi = (1 + \alpha^2)^{1/2} \frac{(\gamma^2 - 1)^{1/2}}{\gamma + 1} \left(1 + \frac{\Delta}{2} \frac{\alpha^2}{1 + \alpha^2} \frac{\alpha^2 - \gamma}{\gamma} \right) \frac{\omega L}{c} \frac{dV}{V}, \quad (1)$$

where $d\phi$ is the phase fluctuation, $\alpha = v_{\perp}/v_z$ is the perpendicular to the axial velocity ratio, γ is the relativistic energy factor, ω is the signal frequency, L is the gyroklystron circuit length, c is the speed of light, dV/V is the fractional voltage ripple, and

$$\Delta = \frac{\gamma^2}{\gamma^2 - 1} \frac{\alpha^2 + 1}{\alpha^2} \frac{2}{\omega} \left(\omega - \frac{\Omega}{\gamma} \right)$$

where Ω/γ is the electron cyclotron frequency, Δ is called the normalized detuning parameter, and gyrotrons are usually designed so that optimum operation occurs when Δ is in the range between 0 and about 1.

It is clear from Eq. (1) that when comparing two gyroklystrons, say one designed to work at low voltage ($\gamma \approx 1$) with one designed to work at higher voltage, the low voltage gyroklystron may be expected to operate with relatively small values of phase fluctuation. Indeed, we have measured the phase fluctuation[3] in a 30 kV three-cavity gyroklystron located at the Naval Research Laboratory[4], and have found $d\phi = 4^\circ$ for each 1% fluctuation in power supply voltage. On the other hand, for the 500 kV four-cavity gyroklystron which was studied at the University of Maryland under the present contract, numerical calculation[5] yielded $d\phi = 28.5^\circ$ for each 1% ripple in power supply voltage.

Both of these values of phase fluctuation are in good agreement with calculations made using Eq. (1). Thus, if 500 kV gyrotrons are to be employed in driving high energy linacs, feedback to reduce phase fluctuation is of compelling interest.

A fast feedback loop as depicted in Fig. 31 was employed[6] to reduce phase fluctuations in the 30 kV gyroklystron mentioned above. Both proportional and proportional-integral type feedback circuits were used.

In Fig. 32, phase fluctuation of the gyroklystron output signal is plotted vs. feedback gain for the proportional feedback circuit. The feedback reduces the phase fluctuation from

Figure 31: Experimental setup for phase jitter reduction by feedback circuit.

1.9 degrees to 0.9 degrees corresponding to a fluctuation reduction ratio of 0.47. A superior phase fluctuation reduction ratio of 0.3 was obtained by adding an integrating circuit in the feedback loop.

Using these results, one can estimate the degree of power supply ripple which could be tolerated in 500 kV gyrokystron used as linac drivers. If the tolerable rms phase fluctuation was $d\phi = 5^\circ$ and a phase fluctuation reduction ratio of 0.3 were achievable by using feedback, then the tolerable voltage ripple would be

$$\frac{dV}{V} = \frac{5^\circ}{28.5(^\circ/\%) \times 0.3} = 0.6\%.$$

This is a reasonable power supply specification.

Figure 32: Phase fluctuation vs. feedback gain in a gyrokystron.

III. Theoretical Research Progress

A. Nonlinear Code Update

For the last five years our primary design tool was our nonlinear gyrokystron amplifier code — GYKL. During that time we have added various improvements, the most notable being space charge depression, realistic field profiles, and AC space charge[7]. Recently we have obtained experimental data which can be compared to the predictions of GYKL, at least in the linear regime. Figure 33 shows experimental and theoretical results for the two-cavity experiment (tube IV as described in Section II.C) with the following parameters: $V_b = 350$ kV, $I_b = 125$ A, $\alpha = 0.8$, and $B_0 = 5.09$ kG. The dashed line in this figure is the theoretical prediction with AC space charge included and the solid line is the experiment. Agreement is fairly good, especially given the experimental uncertainties (input power, pitch angle, and velocity spread are known only approximately). The almost exact match between experiment and theory at 30 kW input power is probably fortuitous; at this input power we believe that there is a shift in the field profile which enhances the efficiency and GYKL cannot model such a shift at this time. Thus, in the region where we believe our code, the predicted gain is higher than the experimental gain by about a factor of two.

Figure 33: Comparison of theory and experiment for a two cavity gyrokystron amplifier

B. Mode Suppression

In the past year we have seen experimentally that mode competition is one of the greatest obstacles to achieving our design goal of 30 MW. Mode competition, at least in the linear regime, is in principle a straightforward numerical problem: find the fields in the cavity and use those fields to compute the average change in energy of an ensemble of particles as they traverse the cavity. However, because of the complexity of our system, computing the electromagnetic fields is difficult and integrating the particles equations of motion is costly in terms of computer time. Thus, our theoretical work in the last year has concentrated on further development of two codes: CASCADE, which uses the scattering matrix formalism to compute the fields in a complex cavity containing lossy dielectrics[1]; and QPB, which uses the Lie transform method[8] to compute the start oscillation currents for the fields produced by CASCADE. These codes are discussed in more detail below.

1. CASCADE

We were working on CASCADE for the past several years. Until recently, it has suffered from two major drawbacks:

1. It could not handle cavities containing dielectrics with high dielectric strengths.
2. It had not been adequately tested.

In the past year, we have solved item (1). CASCADE can now compute the fields in cavities with dielectric constants as large as 40 in reasonable geometries, certainly large enough for our experimental requirements.

Testing CASCADE was more difficult, although we have made some progress. At present, there are no other codes in existence (at least that we are aware of) that can adequately handle *lossy* dielectrics, so comparison with existing codes is not possible; our code was compared with URMEL-T for purely *real* dielectrics, and agreement was excellent. Thus we were concentrating on experimental verification. So far the agreement between CASCADE and experiments was fair; we were hampered mainly by the difficulty in constructing homogeneous lossy dielectrics and obtaining accurate values for the real and imaginary parts of dielectrics which we do believe are uniform. A typical set of data is shown in Fig. 34. There is certainly approximate agreement, but we do not know at this time whether the discrepancies are caused by problems with the code or uncertainties in the values and/or homogeneity of the dielectrics. At this point our sentiment is that the code is correct, although further comparisons are clearly needed.

Finally, we note that when there are no dielectrics present, the agreement between CASCADE and experiment is essentially perfect, as shown in Fig. 35. We hope to achieve the same kind of agreement with lossy dielectrics present.

2. QPB

In the past, to compute the start oscillation current for a given mode, we would directly integrate the equations of motion for an ensemble of particles. The problem with this approach was that it took about one minute of computer time for every magnetic field point. This was fast enough to characterize existing circuits, but too slow to perform detailed parameter studies or include the effects of velocity spread.

Since we are interested in linear mode competition, it is in principle possible to use perturbation theory to compute the start oscillation current analytically. However, conventional

Figure 34: Comparison of CASCADE with experiment for a complex cavity consisting of six lossy dielectric rings (inner radius = 1.5 cm, outer radius = 3.1 cm, $\epsilon = 9.2 + 0.4 i$, length = 0.74 cm) separated by five sections of metal waveguide (radius = 1.5 cm, length = 0.74 cm).

Figure 35: Comparison of CASCADE with experiment for a complex cavity without lossy dielectrics. This plot corresponds to Fig. 7 of Ref. 1.

Figure 36: Schematic of the type of complex cavities we model.

perturbation methods are simply intractable in complex cavities of the type we employ. To get around this problem we use Lie transform perturbation theory, which makes explicit use of the Hamiltonian structure of the equations of motion. With this method, we were actually able to derive an *analytical* expression for the start oscillation current in complex, cylindrical, dielectric loaded cavities. Figure 36 shows a schematic of the kinds of cavities we model. Each section, labelled by ℓ , is cylindrically symmetric and contains an outer layer of dielectric with arbitrary thickness and dielectric constant. To compute the fields in the whole cavity, we use the scattering matrix formalism:[1] the fields are written as an eigenmode expansion in each section, and the mode amplitudes are determined by enforcing continuity of the fields across a section boundary. At the ends, we take outgoing boundary conditions. In a given section ℓ , the z -components of the electric and magnetic fields are written

$$E_z = \text{Re} \left[-ie^{-i\omega t} \sum_n \left(f_n e^{ik_n(z-z_L)} - b_n e^{-ik_n(z-z_R)} \right) C_{En} J_m(k_{\perp n} r) e^{im\theta} \right]$$

$$H_z = \text{Re} \left[-ie^{-i\omega t} \sum_n \left(f_n e^{ik_n(z-z_L)} + b_n e^{-ik_n(z-z_R)} \right) C_{Hn} J_m(k_{\perp n} r) e^{im\theta} \right]$$

where ω is the real part of the annular frequency, f_n and b_n are the scattering amplitudes computed by CASCADE, C_{En} and C_{Hn} are the mode amplitudes for TM and TE modes, respectively, J_m is the Bessel function of order m , m is the azimuthal mode num-

ber, $k_{\perp n}$ is the perpendicular wavenumber for radial mode n , k_n is the z -component of the wavenumber, and z_L and z_R are the left- and right-most positions of section ℓ . The quantities f_n , b_n , C_{En} , C_{Hn} , $k_{\perp n}$, k_n , z_L , and z_R should really have an additional subscript ℓ denoting section number; we drop that subscript for clarity.

The start oscillation current, $I_{s.o.}$, for this type of cavity is determined from energy balance; its definition is

$$I_{s.o.} \equiv \frac{\omega W_{Em}}{Q V_b \eta}$$

where W_{Em} is the electromagnetic field energy, Q is the cold cavity quality factor, V_b is the beam voltage, and the efficiency, η , is defined by

$$\eta \equiv \frac{\overline{\Delta\gamma}}{\gamma - 1}.$$

The quantity γ is the initial relativistic factor (i.e., $mc^2(\gamma - 1)$ is the initial particle kinetic energy), and $mc^2\overline{\Delta\gamma}$ is the difference between the initial and final kinetic energy averaged over the particles. The significance of the start oscillation current is that for beam currents greater than $I_{s.o.}$, the mode with frequency ω and quality factor Q is unstable, while for beam currents below $I_{s.o.}$, the mode is stable.

The energy shift $\overline{\Delta\gamma}$, which is second order in the amplitude of the electromagnetic field, can be computed using the Lie transform method. The result, which is actually not much more complicated than the one for the right circular cylindrical cavity, is written

$$\overline{\Delta\gamma} = -\frac{\omega}{4P_0} \left(\frac{c}{\omega}\right)^3 \sum_s \left[\partial_\gamma - \frac{2sc^2}{\omega\Omega_0 r_L} \partial_{r_L} + \frac{2(m-s)c^2}{\omega\Omega_0 r_g} \partial_{r_g} \right] \left| \sum_{n,\ell} \Delta W_{snt} \right|^2$$

$$\Delta W_{snt} = f_n \chi(k_n, C_{En}, C_{Hn}) + b_n e^{ik_n L} \chi(-k_n, -C_{En}, C_{Hn})$$

$$\begin{aligned} \chi(k_n, C_{En}, C_{Hn}) &= \frac{\omega v_z/c}{\omega + s\Omega_c - k_n v_z} e^{-i(\omega + s\Omega_c)z/v_z} \\ &\times \left[e^{-ik_n L - i(\omega + s\Omega_c)L/2v_z} - e^{i(\omega + s\Omega_c)L/2v_z} \right] \\ &\times J_{m-s}(k_{\perp n} r_g) \left[C_{Hn} \frac{\Omega_c}{k_{\perp n} v_z} \frac{\omega r_L}{c} J'_s(k_{\perp n} r_L) - C_{En} \left(\frac{s\Omega_c}{k_{\perp n} v_z} \frac{k_n}{k_{\perp n}} + 1 \right) J_s(k_{\perp n} r_L) \right]. \end{aligned}$$

In this expression, $P_0 = m_e^2 c^5 / q^2 = 8710$ MW, c is the speed of light in vacuum, $\Omega_0 = qB_0/m_e c$ is the nonrelativistic cyclotron frequency, B_0 is the applied magnetic field, $\Omega_c = \Omega_0/\gamma$ is the relativistic cyclotron frequency, r_L is the Larmor radius, r_g is the guiding center

radius, v_z is the velocity, L is the length of section ℓ , and $\bar{z} = (z_L + z_R)/2$ for section ℓ . For clarity we have suppressed the subscript ℓ , which labels the section number; that subscript should appear on the variables k_n , $k_{\perp n}$, f_n , b_n , C_{En} , C_{Hn} , L , and $\bar{z}$.

Using this expression, we plot the start oscillation current versus magnetic field for tube IV (see Section II.C) including the downtaper region and the output window. We look at eight modes ranging in frequency from 6.09 to 7.47 GHz; all with azimuthal mode number $m = 1$ operating in the first harmonic. The results are given in Figs. 37a–37c. Figure 37a and 37b correspond to a beam voltage of 500 kV and a pitch angle, $v_{\perp}/v_z$, of 1.5; Fig. 37a has no velocity spread while Fig. 37b has a velocity spread in v_z of 10%. Figure 37c corresponds to a beam voltage of 350 kV, a pitch angle of 0.8, and no velocity spread. In all three graphs, frequency is in GHz.

We note that the total Cray time used to compute these graphs was about 1 hour versus an estimated time (using the old method of integrating particles) of about 50 hours.

C. Linear Collider and Gyroklystron Studies (SAIC Collaboration)

During the past seven years, SAIC has played a role in gyroklystron development and linear collider scaling for the University of Maryland. This summary will describe the major progress made over the past 18 months.

1. Linear collider scaling studies

Detailed computer models of the scaling of rf linear colliders were developed to assess the role that a gyroklystron driver might play in a future multi-TeV collider. These models were used to examine the scaling of the linear collider in energy, luminosity, and rf frequency. An interesting early result of the study was the realization that high accelerating gradient (in excess of 100 MV/m) would be counter-productive because the required rf power scales as the square as the gradient while the energy gain is only linear in the gradient. This scaling can be outweighed by the difficulties associated with controlling the wake fields in a very long, low-gradient accelerator. In fact, the SLC and the projected NLC are more susceptible to wake field effects than later, higher gradient colliders will be.

A major thrust of these studies was to examine the frequency scaling of the collider. As

(a)

(b)

(c)

(d)

Figure 37: Start oscillation current in A vs. magnetic field in kG. All frequencies are in GHz.

a) Beam parameters are $V_b = 500 \text{ kV}$, $v_{\perp}/v_z = 1.5$, $\Delta v_z/v_z = 0$; b) beam parameters are $V_b = 500 \text{ kV}$, $v_{\perp}/v_z = 1.5$, $\Delta v_z/v_z = 0.1$; and c) beam parameters are $V_b = 350 \text{ kV}$, $v_{\perp}/v_z = 0.8$, $\Delta v_z/v_z = 0$. d) Tube IV used to compute s.o. curves for graphs a)-c).

one progresses from S band to X band to K band, the device becomes more energy efficient, but it also becomes more sensitive to beam loading effects generally, and also more difficult to fabricate and align.

The most recent study has focused on the scaling from the NLC (0.5 TeV in the center-of mass) to an NLC upgrade (1.5 TeV (CM)) and onward to a 3 TeV (CM) collider, with the system luminosity scaled as the square of the energy to preserve the counting rate. The sensitivity to wake fields at the low gradient of the NLC (33 MV/m) dictates that the frequency not exceed X band (≈ 12 GHz). The NLC Upgrade, with ≈ 100 MV/m gradient, however, can operate up to K_i band (≈ 20 GHz), provided that the SLAC structure was modified to reduce high-order wake fields responsible for the multi-bunch beam-breakup instability. Furthermore, building the NLC Upgrade at K_i band will make it expandable to a 3 TeV (CM) collider. This last device requires an efficiency obtainable only through high-frequency operation, and therefore should be designed at ≈ 20 GHz.

The character of the studies made was to reduce the parameter space to a plane, using known relationship together with the specification of known quantities, and to plot curves representing "constraints" on the parameter plane. For example, the eleven quantities which describe the interaction point are reduced to seven independent quantities by virtue of the relationships that define the luminosity, the disruption parameter, the beamstrahlung loss, and the average beam power. Specifying five of these quantities then reduces the problem to a two-dimensional plane; in SAIC's studies the $\sigma_y - \sigma_z$ plane (transverse vs. longitudinal bunch size) was utilized. The constraint inequalities map to allowed and forbidden regions of the plane. A point in the allowed region is chosen for the interaction point, and the study proceeds to an accelerator model. Here, a convenient plane turns out to be the $E_0 - \lambda$ plane (peak accelerating electric field vs. rf wavelength). Constraint curves representing breakdown, maximum practical levels of average ac power and peak rf power are plotted to this plane together with curves that describe longitudinal and transverse wake-field effects (energy spread and transverse emittance growth). Again the allowed region of the plane will correspond to accelerators which meet all of the constraints. Typically, the allowed region is closed at high frequencies by the transverse wake-field effects and is closed at low frequencies by the ac or rf power constraints. Since the wake-field limit is a very sharp cutoff, one gains

by operating at high frequencies.

These results have consistently indicated that the optimal operating regime for the linear collider is ≈ 20 GHz, where wake-field effects are still tolerable, and the collider will have the efficiency needed to reach multi-TeV energy at high luminosity. The gyrokystron source offers the capability to reach such frequencies (which are likely to be inaccessible to klystrons).

2. Linear collider cost analyses

A cost analysis for the rf drivers and their associated pulse compression systems only was presented at the SLAC Workshop on the Next Generation Linear Collider (December, 1988). That analysis was expanded to include cost studies for the NLC, the NLC Upgrade, and the 3 TeV (CM) collider. This study verifies SLAC's claim that the NLC is designed close to the economic optimum where the length-associated costs are equal to the rf power-associated costs. It also provides several interesting conclusions about the strategy needed to achieve the 3 TeV collider at a reasonable cost. Since a K_i band NLC Upgrade is expandable to the 3 TeV collider, and since it costs no more to build the NLC Upgrade at K_i band than it does at X band, the savings attained by building the # TeV collider as an expansion of an NLC Upgrade built with 20 GHz gyrokystrons rather than building it from an NLC Upgrade built with 12 GHz klystrons will be well in excess of one billion dollars.

This option exists only if a 100 MW peak power gyrokystron at 20 GHz can be developed and demonstrated to have reasonable lifetime and cost by the time the NLC Upgrade is ready for final design.

3. Effect of energy spread on the beam breakup instability

The use of Landau damping (or more properly, BNS damping) was recognized as the means of controlling the single-bunch beam breakup instability in a linear collider. The damping effect occurs when an energy spread is deliberately introduced on the bunch so that the tail has a lower energy than the head. The negative chromaticity of the linac lattice then assures that the tail will be more strongly focused by the applied focusing fields. Since the transverse wake fields tend to defocus the tail more than the head, these two effects can be

made to approximately cancel each other.

Early theoretical studies of this damping mechanism included simulations by Bane and by Henke, and analytical studies based on a two-particle model for the bunch (where the bunch is modeled as two point particles, one at the head and the other at the tail), and more serious theoretical models by Chao, Richter and Yao, and by Balakin, Novokhatsky and Smirnov. The Chao-Richter-Yao model did not include energy spread, while the Balakin-Novokhatsky-Smirnov model assumed the transverse wake field to be constant over the bunch (the actual transverse wake field increase linearly from zero at the beam head).

The calculation carried out at SAIC includes energy spread, a linearly-varying wake field, and acceleration of the bunch in an analytical mode. The model shows that the oscillation amplitude of the bunch stops growing after a finite distance that depends on both that magnitude and the sign of the energy spread. An asymptotic expansion of the exact result is obtained that is valid for large propagation distances and yields an approximate criterion for the validity of BNS damping.

4. Gyrokystron cold-test calculations in three dimensions

The large, overmoded cavities used in gyrokystrons enables them to handle large power at high frequency, but introduces the problem of eliminating spurious oscillations. The experimental program at the University of Maryland is dealing with this issue by introducing lossy elements in the wall of the cavity to selectively damp unwanted modes. SAIC has attempted to assist this portion of the experimental effort by running three-dimensional simulations with the ARGUS code to benchmark the results of an in-house design code at the University of Maryland. The results from ARGUS indicate that the lossy elements planned for the experiment should eliminate the mode completely. Test calculations on model problems showed reasonable agreement between ARGUS and the University's internal code, which will be used for future designs.

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