Quasi-Optical Power Combining of Solid-State Millimeter-Wave Sources

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Abstract —**Very efficient power combining of solid-state millimeter-wave sources may be obtained through the application of quasi-optical resonators and monolithic source arrays. Through the theory of reiterative wavebeams (beam modes) with application of the Lorentz reciprocity theorem, it is shown that planar source arrays containing 25 individual elements or more result in very efficient power transfer of energy from the source arrays to the fundamental wave-beam mode. It is further shown that for identical sources within a properly designed quasi-optical power combiner, the output power tends to increase much faster that number of source elements.**

I. Introduction

CONVENTIONAL waveguide power combiners are limited in power output, efficiency, and number of sources that may be combined in the millimeter-wave region. This limitation is a consequence of the requirement that linear dimensions of conventional waveguide resonators be of the order of one wavelength to achieve acceptable mode separation and to avoid multimode operation. On the other hand, quasi-optical resonators have linear dimensions large compared to wavelength and they offer an attractive approach to overcome these limitations. Fundamental limitations of power combining utilizing quasi-optic resonator techniques is discussed in this paper, and it is shown that very high combining efficiency may be obtained. The approach utilizes an array of source elements placed within a transverse plane near one reflecting surface of the resonator. Energy is extracted from the system through one reflector which is partially transparent.

II. Combiner Configuration

To investigate the feasibility of quasi-optical power combining of millimeter-wave sources, an approach which combines a wavebeam resonator or Fabry–Perot resonator is used as the combining element and sources are modeled as an array of current elements within the resonant structure, as shown in Fig. 1. A wave-beam resonator of rectangular symmetry is utilized and power is extracted from the source array to the lowest order or "Gaussian" mode of the resonator. The resonator consists of two surfaces which are large in terms of the operating wavelength. One surface is a perfect, planar reflector and is located in the plane $z=0$; the other reflector, located at $z=D$, is partially transparent and curved. Useful energy will "leak" through this reflector with a well-defined spatial distribution. The reflector curvature may be expressed by a pair of focal lengths which define the curvature in two perpendicular axial planes, usually the $x-z$ and the $y-z$ planes. The sources are placed in a transverse plane between the reflectors and slightly displaced from the plane reflector. It is assumed that each source, which may be an IMPATT or GUNN diode, is attached to a short dipole which also lies in a transverse plane. A planar array of source diodes with connecting dipoles lends itself to integrated-circuit fabrication techniques [1]. Feedback coupling or signal interaction occurs between the resonant mode and the individual sources leading to injection locking and single-frequency operation. The coupling coefficient of the source array for each mode is calculated through application of the Lorentz reciprocity theorem. Also, the driving point resistance of each dipole in the presence of all other excited dipoles is calculated.

For this configuration, one must consider the electromagnetic fields within two regions of space. Between the reflectors, $0<z<D$, a resonant field exists which consists of two traveling waves, one propagating in the $+z$ or "forward" direction and a second equal amplitude wave traveling in the $-z$ or "backward" direction. The sum of the traveling waves may be expressed as a standing wave whose transverse distribution is described as a sum of the "wavebeam modes." In the region $z>D$, only waves traveling in the $+z$ direction exist, and contain the same spectrum of modes as the fields within the resonator.

III. Theory

A. Electromagnetic Wavebeams and Resonators

Quasi-optic resonators are based upon reiterative wave beams or beam modes. These modes were first described by Goubau and Schwering [2] and they satisfy orthogonality relations like the wave modes in conventional tubular waveguides. In directions transverse to the direction of propagation, characteristic dimensions of fields contained within wave-beam resonators are much larger than those in conventional waveguides. They range from about 20 to many thousand wavelengths depending on the frequency and structures used. In the millimeter/sub-millimeter range, the transverse dimensions are typically from 20 to 100 wavelengths.

Modes of rectangular symmetry are utilized for this investigation since the beam modes, as well as source

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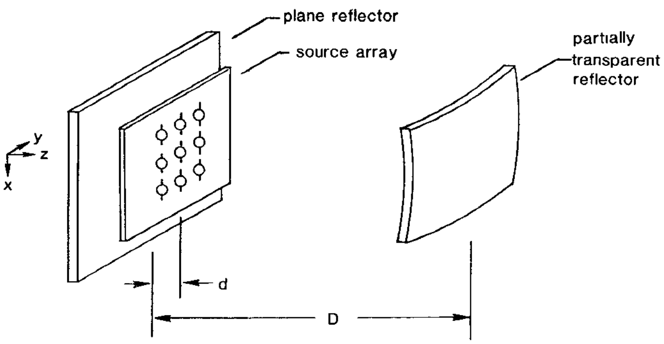

Fig. 1. Resonator–source array configuration.

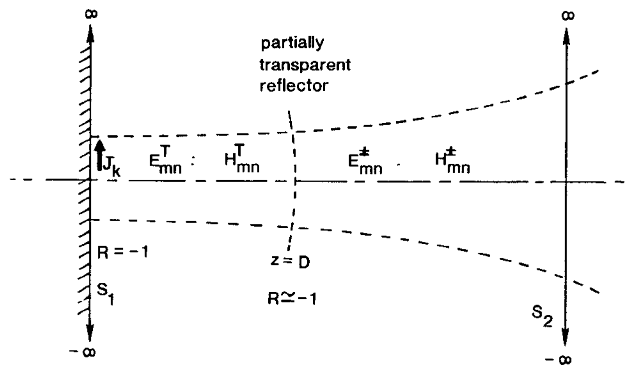

Fig. 2. Cross section showing spatial regions.

coordinates of a regular rectangular array, may be expressed in Cartesian coordinates. Wave beams expressed in Cartesian coordinates are satisfied by Hermite–Gaussian functions [3], [4]. Since the definition of the Hermite polynomials is not uniform in the literature, the following definition is used [3], [5]:

$$H_{en}(X) = (-1)^n (X^2/2) \frac{d^n}{dX^n}(\exp(-X^2/2)). \quad (1)$$

The following recurrence relation is also useful:

$$H_{e(n+1)}(X) = XH_{e(n)}(X) - nH_{e(n-1)}(X). \quad (2)$$

The Hermite polynomials form a complete system of orthogonal functions within the range $-\infty \leqslant X \leqslant \infty$ with the weight function $\exp(-X^2/2)$. An ortho-normal spectrum of wave-beam modes may be obtained from this definition, and is shown below for each linearly polarized component of the wave beam [4]*:

$$\begin{aligned} E^{\pm}_{mn}(x, y, z) &= \frac{(\mu/\epsilon)^{1/4}}{\sqrt{\pi \bar{X}\bar{Y} m! n!}} (1+u^2)^{-1/4}(1+v^2)^{-1/4} \\ &\cdot H_{em}(\sqrt{2}\, x/x_z) H_{en}(\sqrt{2}\, y/y_z) \\ &\cdot \exp\Bigg\{ -\frac{1}{2}\left[(x/x_z)^2 + (y/y_z)^2\right] \\ &\mp j\Bigg[kz + \frac{1}{2}\left(u(x/x_z)^2 + v(y/y_z)^2\right. \\ &-\left(m+\frac{1}{2}\right)\tan^{-1}(u) - \left(n+\frac{1}{2}\right)\tan^{-1}(v)\Bigg]\Bigg\} \end{aligned} \quad (3)$$

where

$$u = z/k\bar{X}^2, \quad v = z/k\bar{Y}^2$$

$$x_z = \bar{X}^2\left(1+\frac{z^2}{k^2\bar{X}^4}\right), \quad y_z = \bar{Y}^2\left(1+\frac{z^2}{k^2\bar{Y}^4}\right)$$

and the relationship between field components is

$$E^{\pm}_{xmn} = \pm\sqrt{\frac{\mu}{\epsilon}}\, H^{\pm}_{ymn}, \quad E^{\pm}_{y} = \mp\sqrt{\frac{\mu}{\epsilon}}\, H^{\pm}_{xmn}. \quad (4)$$

*The argument (x, y, z) of beam modes will be suppressed throughout this paper except when it is necessary to refer to a specific point, such as the location of a current source.

The $E^{\pm}_{mn}$ fields represent the desired wave-beam modes and the + sign refers to traveling waves progressing in the positive z direction, and the − sign refers to waves traveling in the negative z direction. The subscript x or y refers to the polarization. Quantities $\bar{X}$ and $\bar{Y}$ which determine the decay of the field in the x and y directions are called mode parameters. Mode parameters are parameters which are adjusted so that the wavebeam satisfies an imposed condition. When one considers a resonator structure, the condition that must be satisfied is that for each round trip of a wave within the resonator, the field repeats itself in both phase and amplitude distribution. It has been shown that the mode parameter is a function of resonator configuration and wavelength. For the resonator described above, the mode parameters are [4]

$$k\bar{X}^2 = \sqrt{(2 - D/F_x)F_x D} \quad (5)$$

$$k\bar{Y}^2 = \sqrt{(2 - D/F_y)F_y D} \quad (6)$$

where

$k = 2\pi/\lambda$
D = distance between the reflecting surfaces,
F_x = focal length of the curved reflector referenced to the x axis,
F_y = focal length of the curved reflector referenced to the y axis.

The modes satisfy, in any plane z = constant, the orthogonality relations

$$\sqrt{\frac{\epsilon}{\mu}} \iint_{-\infty}^{\infty} E_{mn} \cdot E^{*}_{m'n'}\, dx\, dy = \delta_{mm'}\delta_{nn'}. \quad (7)$$

Since the Hermite–Gaussian functions form a complete system of orthogonal functions, any beam whose transverse electric field is known in a plane z = constant can be expanded into a series of wave-beam modes, providing the beam satisfies a paraxial propagation condition. Experience, however, has shown that this requirement is satisfied in practical systems and that the error tends to occur in higher order modes which usually are not of interest.

The modes described by (3) represent waves traveling freely in space. With reference to Fig. 2, they describe the fields outside of the resonator, or in the region $z > D$. In order to satisfy the boundary conditions within the resonator region, $0 \leqslant z \leqslant D$, one must take for each mode a sum of "forward" and "backward" traveling waves. Under resonant conditions, the fields within the resonator may build up due to multiple reflections, and the amplitude will be limited by the reflection coefficient of the partially transparent reflector. For application of the Lorentz reciprocity theorem that will follow, it is required to determine the resonator fields when excited by a properly adjusted wave beam consisting of a single mode or spectrum of modes incident from $z=\infty$ upon the resonator. These fields become the test fields. In order to compute the worst-case fractional power coefficient, it is assumed that all modes resonate simultaneously. Because the total phase shift of a wave traveling from one reflector to the other depends upon the mode numbers m and n [4], all modes will not resonate simultaneously in an actual wave-beam resonator. The condition of simultaneous resonance is necessary to determine the best location of source elements. This formulation results in the worst-case fractional power-coupling coefficient. The fractional power-coupling coefficient is defined as the power in the desired mode, usually the fundamental Gaussian mode, divided by the power in all modes excited by the given current distribution.

The partially transparent reflector must be characterized in order to determine the electromagnetic field within the resonator. This reflector may be considered as a lossless two-port junction. The scattering matrix of such a junction has certain well-defined properties listed below [6]:

$$S_{11}=S_{22}=Re^{j\theta} \tag{8}$$

$$S_{12}=S_{21}=\sqrt{1-R^2}\,e^{j(\theta+\pi/2)}. \tag{9}$$

Now, it is postulated that for a wave beam incident upon the resonator from $z=\infty$, the wave beams within the resonator have both amplitude and phase differences from the incident wave beam. Since a perfect reflector is located at $z=0$, there is no net power flow through any plane $z>0$. Using the condition of zero net power flow through any transverse plane along with the properties of the lossless partially transparent reflector, the field within the resonator becomes

$$E_{mn}^T = A\,\mathrm{Re}(E_{mn}^+)\sin(kz) \tag{10}$$

where

$$A=\frac{2R\sin(\psi)+\sqrt{1-R^2\cos^2(\psi)}}{\sqrt{1-R^2}}\,e^{j(\psi+\pi/2)}. \tag{11}$$

As seen by (11), the field strength is at its maximum when $\psi=\pi/2$. For this value of ψ, the fields within the resonator are real; thus, the system is considered to be resonant.

B. Coupling to an Array of Current Sources

One can now determine the coupling coefficient to a current element or to an arbitrary array of current elements through application of the Lorentz reciprocity theorem [7] with the further requirement that all current sources are located within the resonator. There is no loss in generality by considering only modes where $E_{ymn}=H_{xmn}=0$, and the impressed currents are x-directed

$$\oint_S \left(E_{mn}^{\pm}\times H_1 - E_1\times H_{mn}^{\pm}\right)\cdot n\,da = \iiint_{V_R} J\cdot E_{mn}^T\,dv \tag{12}$$

where

$E_{mn}^{\pm}$ and $H_{mn}^{\pm}$ = modal fields in space,
E_{mn}^T and H_{mn}^T = fields within the resonator expressed in terms of the modal fields,
E_1 and H_1 = fields in space due to the current elements,
V_R = volume bounded by the resonator.

The method used to find the field radiated by an arbitrary array of filamentary currents within a quasi-optical resonator is to expand the radiated field in terms of normal beam waveguide modes (Hermite–Gaussian functions) and to determine the amplitude coefficients in this expansion. With reference to Fig. 2, let J_k represent an arbitrary infinitely thin current element. Such a current must be maintained by some external source (e.g., an IMPATT or GUNN diode), but in the evaluation of the coupling to beam modes, only radiated fields are of interest, and, consequently, the source which maintains the specified current does not enter the picture here.

The field radiated in the positive z direction by the array of x-directed current elements may be represented by

$$E_1=\sum_{kq} a_{kq}E_{kq}^+\hat{x} \qquad \text{for } z>D \tag{13}$$

$$H_1=\sqrt{\frac{\epsilon}{\mu}}\sum_{kq} a_{kq}E_{kq}^+\hat{y} \qquad \text{for } z>D. \tag{14}$$

Since there is a perfectly conducting plane located at $z=0$ as shown in Fig. 2

$$E_{mn}^T(x,y,\phi)=0. \tag{15}$$

The volume chosen over which it is required to evaluate the Lorentz reciprocity relation is bounded by a surface S which extends to infinity in the transverse directions and consists of an infinite, perfectly conducting plane S_1, located at $z=0$ and a second infinite plane S_2, located in some plane $z>D$. When one then performs the integration over this "closed surface", there is only a contribution by the integrals evaluated on S_2. There is no contribution to the integral over S_1 since the $n\times E=0$ along that surface.

$$\oiint_{S_2}\left[E_{mn}^{\pm}\hat{x}\times\sqrt{\frac{\epsilon}{\mu}}\sum_{kq}a_{kq}E_{kq}^+\hat{y} - \sum_{kq}a_{kq}E_{kq}^+\hat{x}\times\sqrt{\frac{\epsilon}{\mu}}E_{mn}^{\pm}\hat{y}\right]\cdot\hat{n}\,da = \iiint_{V_R} J\cdot E_{mn}^T\,dv. \tag{16}$$

Since $E_{mn}^- = E_{mn}^{+*}$, one can utilize the orthogonality relation (eq. (7)) for wave beams and perform the integration term by term. Therefore

$$\oiint_{S_2} \left(E_{mn}^{\pm} \times H_1 - E_1 \times H_{mn}^{\pm} \right) \cdot n\, da = 2 a_{mn}. \tag{17}$$

Hence

$$a_{mn} = \frac{1}{2} \iiint_{V_R} J \cdot E_{mn}^T \, dv. \tag{18}$$

Again, if one considers the case where the array consists of an array of filamentary currents, that the currents are all aligned with the electric field, and that the length of each current elements is small compared to the mode parameter, this equation can be written as follows:

$$a_{mn} \simeq \frac{1}{2} \sum_p I_p \Delta X_p E_{mn}^T (x_p, x_p, z_p) \tag{19}$$

where

I_p = the current into the "terminals" of the pth current element,

$E_{mn}^T(x_p, y_p, z_p)$ = the electric field strength of the m, n mode at the location of the pth current element,

ΔX_p = effective length of the pth current element.

Hence

$$\Delta X_p = \frac{1}{I_p} \int I_p(l) \cdot dl_p. \tag{20}$$

Now with the knowledge of the expansion coefficients a_{mn} given by (19) and (10), which relates fields internal to the resonator to the external fields, one can determine the total electric and magnetic fields E_1, H_1 due to an array of current elements.

C. Driving Point Resistance of Each Element

Since the goal is to obtain a technique for efficient power transfer from an array of sources, one must know the driving point resistance to each element and then to match the source to that resistance. It is assumed that the resonator is adjusted for resonance; hence, the reactive component is zero or at least very small. Since the dipole elements will be surrounded by a strong electric field due to resonator, the self impedance of the dipole is neglected. The input impedance of a dipole element in the presence of an electric field (created by all sources) may be expressed as [8]

$$Z_p^T = \frac{1}{I_p^2} \iiint_{V_R} J_p \cdot E^T dv \tag{21}$$

where Z_p^T is the driving point impedance for the pth current element.

A more useful result is the driving-point impedance for a given mode. It has been shown theoretically and verified experimentally that a wave-beam resonator may be adjusted so that only one mode may exist for a given frequency (for example, the mode patterns of lasers [9]). One can, therefore, express the driving-point resistance for each mode as follows:

$$Z_{pmn} = \frac{1}{I_p^2} \iiint_{V_R} J_p \cdot E_{mn}^T \, dv. \tag{22}$$

Again, considering the case of small dipoles of equal equivalent length, the following expression is obtained:

$$Z_{pmn} = 2A(\Delta X)^2 \sin^2(kz_p) \cdot \mathrm{Re}\left[E_{mn}^{+}(x_p, y_p) \right] \sum_q \frac{I_q}{I_p} \mathrm{Re}\left[E_{mn}^{+}(x_q, y_q) \right]. \tag{23}$$

This result also may be obtained through considerations of energy conservation. The power flowing into a dipole element may be represented as the square of its terminal current, multiplied by its driving-point resistance. Now, total power into the system is the sum of the power flowing into all individual elements. When this total power is equated to the power flux of the forward-traveling wave beam, one obtains the same result as shown by (23).

IV. Computed Results

The theory developed above enables one to determine the number of current elements required to obtain efficient transfer of power to any wave-beam mode. Of primary interest is the current source locations within the resonator, their amplitudes, and the driving-point resistance for each element when the lowest order "Gaussian wave beam" is efficiently excited. In this section, two specific cases will be considered. First, the case where all current elements are assumed to have equal current moment, and second, where the current moment amplitude is adjusted such that it is proportional to the field strength of the fundamental mode at its location.

To obtain efficient coupling, the current elements must be distributed in a transverse plane in such a way that power is efficiently transferred to the lowest order mode and very little power is transferred to any of the other modes. The efficiency of coupling may be calculated for a given distribution of currents by computing the power radiated by the lowest order mode and comparing it to the total power radiated. From (13), it is seen that the amplitude of each mode is represented by the coefficient a_{qk}; thus, using (4), the power of each mode may be calculated as follows:

$$P_{qk} = a_{qk} a_{qk}^* \sqrt{\frac{\epsilon}{\mu}} \iint_{-\infty}^{\infty} E_{qk} \cdot E_{qk}^* \, dx\, dy = a_{qk} a_{qk}^*. \tag{24}$$

Since the modes are orthogonal for a given array of current elements, the fractional power of the fundamental mode ($m = 0$, $n = 0$) compared to the total power of all modes becomes

$$FP_{00} = \frac{a_{00} a_{00}^*}{\sum_{qk} a_{qk} a_{qk}^*}. \tag{25}$$

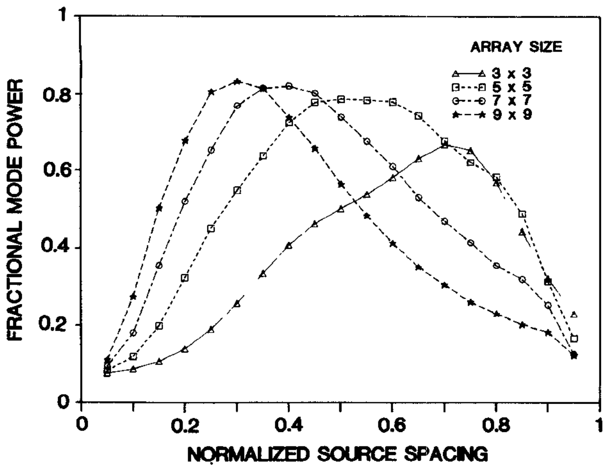

Fig. 3. Fractional power into fundamental mode by equal weight sources.

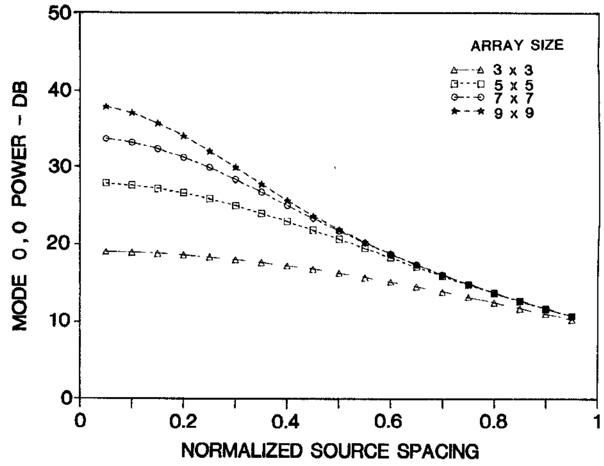

Fig. 4. Power into fundamental mode by equal weight sources.

It should be noted that the excitation coefficient for any mode is determined only by the current distribution and the modal spectrum as if it were freely propagating in space. This restriction is required to obtain the optimum current distribution for the excitation of any given mode. If it were not applied, the mode spectrum would not be complete from the mathematical point of view. From an engineering view point, it represents the worst case since it assumes that all modes are at resonance. Thus, each mode could extract energy from the current elements and is included in the denominator of (25). This assumption clearly aids in determining the optimum source array configuration.

Fig. 3 illustrates the fractional power coupled into the fundamental mode for four different array configurations containing 9, 25, 49, and 81 elements in regular rectangular arrays of equal moment sources with their individual phases adjusted such that each term of (19) is real. All figures that follow have been normalized such that the results presented are independent of the details of the wave-beam resonator; a total of 441 modes are utilized for the computation of the denominator in equation (25). Of course, the normalization must be removed when a particular case is to be considered. To achieve meaningful normalization, the spacing between source elements in each direction is expressed in terms of the wave-beam mode parameter (the $1/e^2$ distance). The source array is considered to lie in a plane transverse to the wave beam and is symmetrical about the wave-beam axis. A practical location for the source array is very close to the reflecting surface located at $z = 0$. For this location, all elements will have uniform phase and the reflecting surface can also become the heat sink for active elements. In terms of coupling energy into the fundamental mode, Fig. 3 shows that for each array configuration there is an optimum source element spacing. It also shows that the maximum source array length for optimum coupling is approximately independent of the number of array elements. The array will extend in each direction from the wave-beam axis about 1.2 mode parameters. Since the ultimate goal is to combine many individual sources to obtain a high power source, the total power

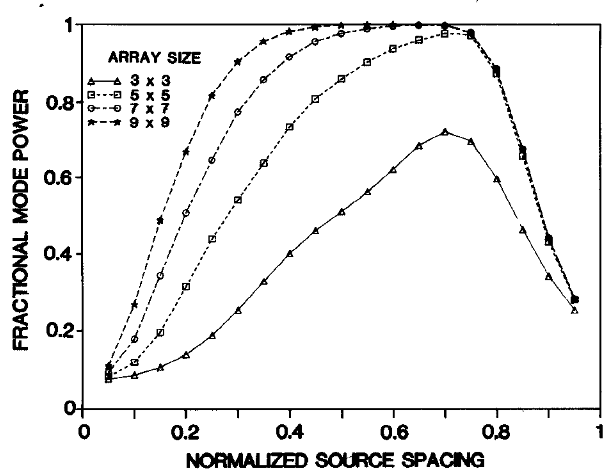

Fig. 5. Fractional power into fundamental mode by Gaussian weight sources.

delivered to the fundamental mode must be determined. Fig. 4 shows the increase of power as the number of sources increases. Zero decibels is the power delivered to the fundamental mode by a single source located on axis. The trend of these curves shows that one should make a tradeoff between array element spacing for optimum fractional power and fundamental mode power. It appears that the source spacing should be reduced so that the optimum fractional power reduces by about 1 dB.

Fig. 5 illustrates the fractional power into the fundamental mode for four different array configurations consisting of regular rectangular source arrays of 9, 25, 49, and 81 elements, and the current moment of each element is adjusted to have a value proportional to the field strength of the fundamental mode at the location of the element (the source array current moments have a Gaussian taper). In this case, very efficient coupling may be obtained since the source array has been matched to the fundamental mode. However, Fig. 6 shows that the fundamental mode power decreases much faster as the source spacing is increased than for the previous case. The net conclusion is that for a power combiner, significant output power reduction will occur if the source spacing is allowed to increase.

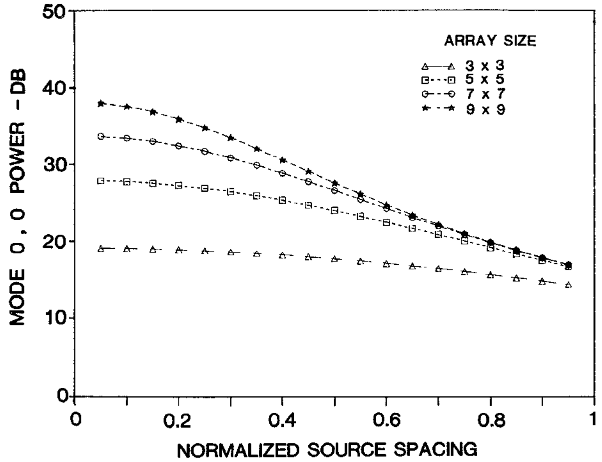

Fig. 6. Power into fundamental mode by Gaussian weight sources.

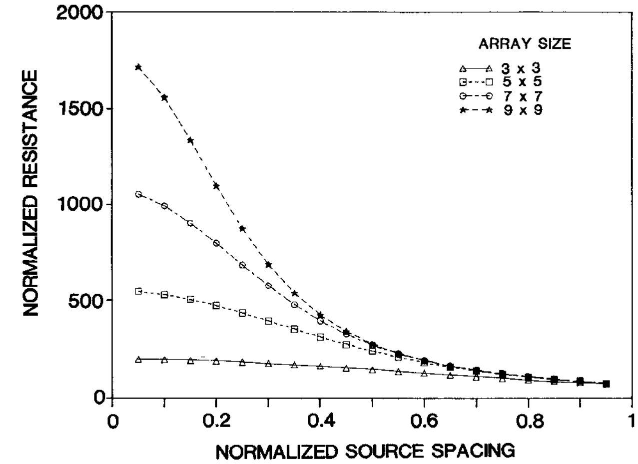

Fig. 8. Driving-point resistances for Gaussian weight sources.

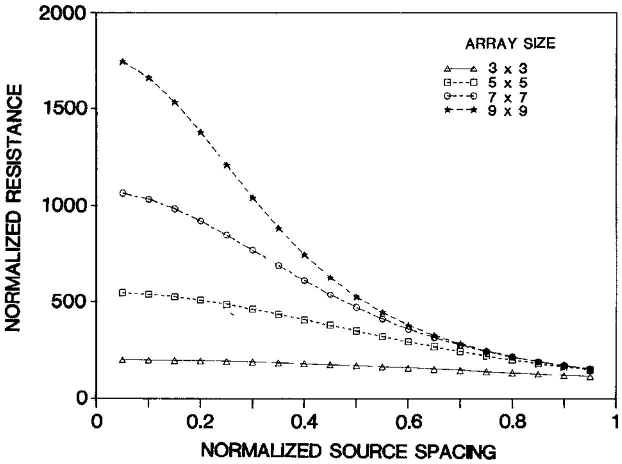

Fig. 7. Driving-point resistances for equal weight sources.

This is especially true for source arrays containing a large number of elements.

Utilizing (23), a family of curves for the driving-point resistance of each element of the source array is obtained. However, this family of curves can be reduced to a single curve for each array configuration. Equation (23) may be written as

$$R_{pmn} = \frac{(\Delta X)^2}{\overline{XY}} \sin^2(kz_p) \sqrt{\frac{1+R}{1-R}} \cdot \left[\mathrm{Re}\left[\sqrt{\overline{XY}}\, E^+_{mn}(x_p, y_p)\right]\overline{R} \right. \quad (26)$$

where

$$\overline{R} = 2\sum_q \frac{I_q}{I_p} \mathrm{Re}\left[\sqrt{\overline{XY}}\, E^+_{mn}(x_q, y_q)\right]. \quad (27)$$

Now (27) represents a normalized resistance factor and depends only upon normalized source spacing, while (26) is the actual driving-point resistance and requires detailed knowledge of the resonator configuration. Fig. 7 is the normalized driving-point resistances for the four array configurations described above, where each element has the same effective length and driving-point current. Fig. 8 is the normalized driving-point resistances for the four array configurations when each element has the same effective length and the terminal current is proportional to the field strength of the fundamental mode at the current element location.

V. Examples

Quasi-optical millimeter-wave power combining was experimentally investigated by Wandinger and Nalbandian [10]. They utilized a wave-beam resonator with two waveguide ports loaded with dielectric rods to couple energy into the system and reported power-combining efficiency of 52 percent. This value is in general agreement with the theory presented here. Each waveguide aperture loaded with a dielectric rod was modeled as four small current elements in a rectangular array separated by 0.1 mode parameters. The location of these "patches" of currents was estimated from the photograph in the paper by Wandinger and Nalbandian to be 0.45 mode parameters from the beam axis. Due to the mode-dependent phase shift of wave-beam modes, only one fourth of the total mode spectrum would simultaneously be resonant in a confocal resonator for a given frequency. All of the above conditions were applied and a coupling efficiency of 40 percent was calculated. Since this theory does not take into account direct, near-field coupling between closely spaced dielectric rod antennas, the agreement is considered good.

Figs. 3 and 5 show that efficient transfer of energy between the array and the wave beam may be obtained for source arrays of a 5×5 and larger if the proper spacing between elements is chosen, while Figs. 4 and 6 show that with the same spacing between array elements there is a diminishing return of power transferred to the fundamental mode as source arrays become larger. The following example is representative. It is assumed that active elements are arranged in the configuration of a uniform 5×5 array and are fabricated as a monolithic structure in GaAs [1]. The transverse dimension of the plane reflector is taken to be 5 cm, which is about the size of available GaAs wafers. The resonator will be "semi-confocal", therefore, $F_x = F_y = D$. The following conditions are also chosen: the

TABLE I
DRIVING-POINT RESISTANCES FOR 5×5 SOURCE ARRAY LOCATED *d* MILLIMETERS FROM PLANE REFLECTOR

R-ohms \ d-mm	0.05	0.1	0.15
R_{00}	2.11	8.43	18.9
R_{10}	1.95	7.78	17.5
R_{11}	1.79	7.78	16.2
R_{20}	1.53	6.12	13.7
R_{21}	1.41	5.65	12.7
R_{22}	1.11	4.44	10.0

mode parameters $\overline{X}$ and $\overline{Y}$ are 1 cm; the operating frequency is 100 GHz; the normalized current element length $\Delta X/\overline{X}$ is 1/50; the normalized spacing between source elements is 0.4; and the reflection coefficient R of the partially transparent reflector is 0.98. From (5) and (6), one obtains $D = 20.9$ cm. The driving-point resistance for each element of the source array is shown in Table I. It should be noted that, because of symmetry, there are only six different driving-point resistances. The array elements all are numbered in matrix notation with the 0,0 element located on the wave-beam axis. For the example shown in Table I, the driving-point resistances were computed by (26) are shown for a 5×5 source array located 0.05, 0.1 and 0.16 mm from and parallel to the plane reflector. In addition, the region of space between the source array and the plane reflector is filled with GaAs. Since IMPATT devices are designed to operate with low driving-point resistances [11], a distance d of 0.1 mm may be chosen as a compromise between the desired low driving-point resistances and the minimum practical thickness of GaAs.

If each active source element is able to maintain the same driving current independent of other nearby sources, and if a single source provides an output power of 1 mw when combined in the quasi-optical power combiner, 25 such sources in a 5×5 array would provide an output power of about 300 mw, 49 such sources in a 7×7 array would provide about 630 mw, and 81 such sources in a 9×9 array would provide less than 800 mw. The above example assumes the separation between source elements remains constant at 0.4 mode parameters and indicates that there may be a diminishing return upon increasing the number of source elements to very large numbers. However, with proper design, one may conclude from this study that it is practical to combine large numbers of millimeter-wave sources using quasi-optical techniques and that substantial power may be obtained.

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