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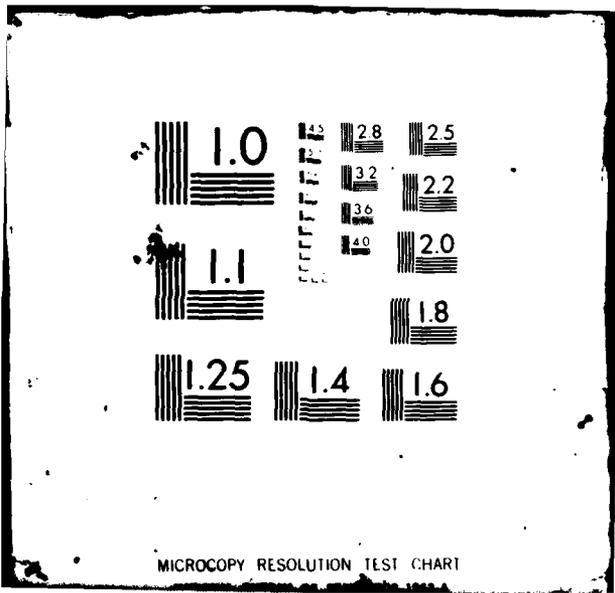
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**NOTE ON PLANAR UNDULATOR WIGGLER
SINGLE ELECTRON MOTION**

BY ROBERT CAWLEY

RESEARCH AND TECHNOLOGY DEPARTMENT

15 JUNE 1981

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axial periodicity of the field, a variable change having some interesting features is reported here. The axial motion between $z = 0$ is converted to the oscillations of a particle in an infinitely long, axially reflecting box. The transit of a half-cycle of wiggler oscillation, $(\Delta z) = \pi/k_0$, corresponds to a single bounce between $Z = \pm \infty$, where Z is the new coordinate. Transverse coordinates and velocities are continuous at each axial reflection from infinity.

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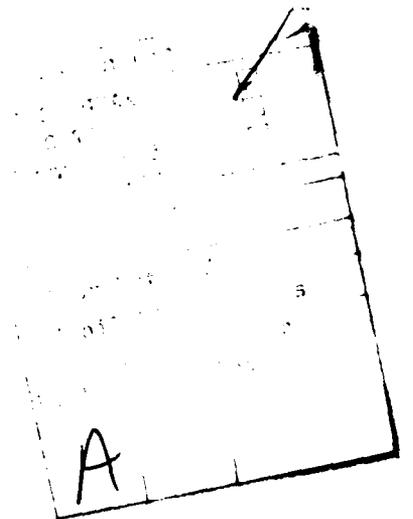
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CHAPTER I

INTRODUCTION

Determination of beam equilibrium configurations forms the starting point in analysis of the free electron laser instability for an intense beam. For equilibrium distributions of beam electrons to be constructed which will take into account nonlinearities associated with beam electron motions in the prescribed magnetic field through which the beam passes, as well as effects of finite transverse geometry, the invariants of single particle motions in the given field are needed. When beam self-fields can be neglected, the determination of such invariants will constitute a solution for the single-particle motion; the equilibrium distribution will be a function of these invariants. Corresponding modulational influences on beam envelope can then be included automatically in the stability analysis.

In two different model field examples this has been done by Bernstein and Hirshfield,¹ Davidson and Uhm^{2,3} and Uhm and Davidson.⁴ In Refs. 1-3 the helical wiggler field was allowed to have only axial dependence,

$$\vec{B}_{\text{wigg}} = -\hat{B} \cos k_0 z \hat{x} - \hat{B} \sin k_0 z \hat{y} \quad (1.1)$$

where $k_0 = 2\pi/\lambda_0$, with λ_0 the wiggler wavelength, and where \hat{B} is the constant amplitude of the wiggler field. In Ref. 3 a uniform axial guide field also was present. In Ref. 4 an axial undulator wiggler with a guide field was studied for the case of a hollow beam, which allowed radial variations of the field components again to be neglected,

$$\vec{B} = B_r \hat{r} + B_z \hat{z}, \quad (1.2)$$

where the axisymmetric components of \vec{B} are got from the $\hat{\theta}$ - component of a vector potential given by $\vec{A} = A_\theta \hat{\theta}$,

$$B_r = -\partial A_\theta / \partial z, \quad B_z = r^{-1} \partial (r A_\theta) / \partial r, \quad (1.3)$$

with $A_\theta = A_\theta(r, z)$, and with also $\nabla \times \vec{B} = 0$. Setting $r = R_0$, the equilibrium beam mean radius, allowed B_r and B_z to be approximated as independent of r , and simplified the analysis of the single particle equations of motion. The field then has a form similar to eq. (1.1),

$$B_r \approx -\hat{B}_r \sin k_0 z, \quad B_z \approx B_0 - \hat{B}_z \cos k_0 z \quad (1.4)$$

where B_0 is the uniform guide field, and \hat{B}_r , \hat{B}_z are constant wiggler amplitudes determined by B_0 , $k_0 R$, and the mirror ratio $R \equiv B_z(r=0, k_0 z = \pm \pi) / B_z(r=0, k_0 z = 0)$.

In Refs. 1 and 2, three exact invariants of electron motion, the energy and two components of transverse canonical momentum, were available; and the same is true of Ref. 3, but where the invariants (C_1 , C_h , C_z), now also allowed incorporation of helical and axial field effects together. In the undulator configuration of Ref. 4, only two exact invariants were available, the energy and canonical axial angular momentum, and a third approximate invariant, the axial component of particle momentum, p_z , was employed. So the undulator wiggler field single particle motions have also proved more difficult.

In the present report I describe a brief investigation of single particle motion in a planar undulator field, specified by a vector potential $\vec{A} = A_y(x, z)\hat{y}$. Only two invariants are immediately available, the energy and \hat{y} -component of canonical momentum. As a starting point of the search for a third invariant with origin in the axial periodicity of the wiggler, or to produce a convenient basis of determination of an approximate third invariant better than p_z , I have transformed the axial coordinate z to a new coordinate Z , which ranges between $+\infty$ as z covers a half-cycle of wiggler oscillation. The transformation is fixed by specifying that the wiggler amplitude be absent from the axial equation (for zero guide field). After elimination of the y coordinate, the xZ -equations describe a particle oscillating in a box of infinite "length" and having axially reflecting walls, while also executing its transverse motions. In this way, the full motion along the beam axis is transformed into motion over only a single cycle of wiggler oscillation.

It is not clear whether the altered problem presents special calculational advantage, and time has not permitted more than a cursory examination of this. It has seemed appropriate nevertheless to record the present formulation as a slightly unusual and possibly

useful starting point of analysis in which features associated with wiggler periodicity receive explicit attention at the outset. The method is more generally applicable, e.g. to problems of cylindrical symmetry as well as to the present case of planar symmetry.

CHAPTER II

PLANAR UNDULATOR WIGGLER EQUATIONS OF MOTION

The vector potential is $\vec{A} = A_y \hat{y}$, and the magnetic field's nonvanishing components are

$$B_x = -\frac{\partial A_y}{\partial z}, \quad B_z = \frac{\partial A_y}{\partial x}, \quad (2.1)$$

The choice

$$A_y(x, z) = B_0 x - k_0^{-1} B_1 \sinh k_0 x \sin k_0 z \quad (2.2)$$

models a static undulator configuration with mean guide field B_0 and undulator amplitude B_1 , where B_0 and B_1 are constants. Due to the symmetry

$$A_y(-x, z) = -A_y(x, z) \quad (2.3)$$

B_z has the sign of x , corresponding to like-pole magnet pairings successively alternating along the \hat{z} - direction. The field satisfies the condition $\nabla \times \vec{B} = 0$. The Hamiltonian,

$$H = c \left[P_x^2 + P_z^2 + \left(P_y - qc^{-1} A_y(x, z) \right)^2 + m^2 c^2 \right]^{1/2} \quad (2.4)$$

where c is the speed of light in vacuo, q the charge on the electron and m its mass, is invariant to translations along the \hat{y} -axis so P_y is a constant of the motion along with H ; and the problem of configuration motion is therewith reduced to a two-dimensional one in x and z .

It will be enough to study the configuration equations of motion. These are

$$\frac{d^2x}{dt^2} = \frac{q}{\gamma mc} B_z(x, z) \dot{y} \quad (2.5)$$

$$\frac{d^2z}{dt^2} = - \frac{q}{\gamma mc} B_x(x, z) \dot{y} \quad (2.6)$$

and

$$\begin{aligned} \frac{d^2y}{dt^2} &= \frac{q}{\gamma mc} (B_x(x, z) \dot{z} - B_z(x, z) \dot{x}) \\ &= - \frac{d}{dt} \left(\frac{q}{\gamma mc} A_y(x, z) \right), \end{aligned} \quad (2.7)$$

the last by the constancy of γ , whose value is H/mc^2 , and by eqs. (2.1). From eq. (2.7)

$$\dot{y} = \frac{P_y}{\gamma m} - \frac{qA_y(x, z)}{\gamma mc}, \quad (2.8)$$

where P_y is the constant of integration; whence, upon substitution into eqs. (2.5) and (2.6) one has

$$\frac{d^2x}{dt^2} = - \frac{\partial \psi}{\partial x}, \quad \frac{d^2z}{dt^2} = - \frac{\partial \psi}{\partial z} \quad (2.9)$$

where

$$\psi = \psi(x, z) = - \frac{qP_y}{\gamma^2 m^2 c} A_y + \frac{1}{2} \left(\frac{qA_y}{\gamma mc} \right)^2. \quad (2.10)$$

Eqs. (2.9) have the first integral

$$\frac{1}{2} (\dot{x}^2 + \dot{z}^2) + \psi(x, z) = C = \text{const.} \quad (2.11)$$

P_y , γ and C satisfy a relationship

$$C = -\frac{1}{2} \left(\frac{P_y}{\gamma m} \right)^2 + \frac{1}{2} \left(1 - \gamma^{-2} \right) c^2, \quad (2.12)$$

so only two of the three constants are independent, and γ can be exchanged for C if we wish.

Note from eq. (2.8), $P_y = 0$ corresponds to $\langle \dot{y} \rangle = 0$ for electron distributions even under $x \rightarrow -x$ owing to eq. (2.3). In this case also, $2C$ has the value v_0^2 , where v_0 is the constant speed of the charge.

I assume $B_0 = 0$ for simplicity now and expand eqs. (2.9) using (2.2) and (2.10),

$$\frac{d^2 x}{dt^2} = -\cosh k_0 x \cdot \operatorname{sink}_0 z \cdot (a + b \sinh k_0 x \cdot \operatorname{sink}_0 z) \quad (2.13)$$

$$\frac{d^2 z}{dt^2} = -\sinh k_0 x \cdot \operatorname{cosk}_0 z \cdot (a + b \sinh k_0 x \cdot \operatorname{sink}_0 z), \quad (2.14)$$

where

$$a \equiv qB_1 / (\gamma m)^2 c \cdot P_y = \frac{\Omega_1}{\gamma m} P_y \quad (2.15)$$

$$b \equiv (qB_1 / \gamma m c)^2 k_0^{-1} = k_0^{-1} \Omega_1^2, \quad (2.16)$$

with $\Omega_1 \equiv qB_1 / \gamma m c$; from eq. (2.11) also

$$\begin{aligned} \dot{x}^2 + \dot{z}^2 + 2 a k_0^{-1} \sinh k_0 x \cdot \operatorname{sink}_0 z \\ + b k_0^{-1} \sinh^2 k_0 x \cdot \sin^2 k_0 z = 2C. \end{aligned} \quad (2.17)$$

CHAPTER III
CHANGE OF VARIABLES

If a new axial variable Z is introduced,

$$Z = k_0^{-1} g(k_0 z), \quad (3.1)$$

independent of t so the Hamiltonian will be unchanged in value, and with g a function at our disposal, we have the identities,

$$\dot{z} \equiv g'(k_0 z)^{-1} \dot{Z} \quad (3.2)$$

and

$$\frac{d^2 z}{dt^2} \equiv g'(k_0 z)^{-1} \left[\ddot{Z} - k_0 g''(k_0 z) \dot{z}^2 \right], \quad (3.3)$$

where primes denote differentiation with respect to the argument. Eq. (2.14) becomes

$$\begin{aligned} \ddot{Z} - k_0 g'' \left[2C - \dot{x}^2 - 2ak_0^{-1} \sinh k_0 x \cdot \sin k_0 z \right. \\ \left. - bk_0^{-1} \sinh^2 k_0 x \cdot \sin^2 k_0 z \right] = -g' \cdot a \sinh k_0 x \cdot \cos k_0 z \\ - g' \cdot b \sinh^2 k_0 x \cdot \sin k_0 z \cos k_0 z. \end{aligned} \quad (3.4)$$

Note the choice $g(\beta) = \beta$, with $\beta \equiv k_0 z$ gives $Z = z$, and (3.4) reduces to (2.14) as it should.

The x -dependence of the coefficients of a , to be called a -terms, is the same on both sides of eq. (3.4); the same is true of the coefficients of b , viz. the b -terms. So the a -terms can be transformed away by a suitable choice of g ; or the b -terms can be

removed instead. (They cannot both be removed by a transformation of z .) The former is accomplished if g satisfies

$$2g''(\beta) \sin\beta = -g'(\beta) \cos\beta \quad (3.5)$$

i.e.

$$g'(\beta) = \pm A / \sqrt{|\sin\beta|}, \quad A > 0. \quad (3.6)$$

The transformation cannot be determined as an elementary function owing to the square root in eq. (3.6), whose cause was the factor 2 in eq. (3.5).

For the b-terms matters are different; the analogue of eq. (3.5) is

$$g''(\beta) \sin^2\beta = -g'(\beta) \sin\beta \cos\beta, \quad (3.7)$$

whence

$$|g'(\beta) \sin\beta| = |A|, \quad (3.8)$$

or, more serendipitously,

$$g'(\beta) = A \cos\beta, \quad (3.9)$$

where A is an arbitrary nonzero constant. Integrating,

$$k_0 z \equiv g(\beta) = A \log |B \tan \frac{1}{2} \beta|, \quad B = \text{const.}, \quad (3.10)$$

where, it will be recalled, $\beta = k_0 z$. B may be assumed positive without loss of generality; I assume further that $B = 1$ as there seems to be no gain resulting from its retention as a free parameter. (Figure 1 shows a sketch of $g(\beta)$ for $A > 0$, and $0 < \beta < \pi$.) Solving eq. (3.10) gives also,

$$\sin\beta = \pm \operatorname{sech} A^{-1} k_0 z, \quad \tan\beta = \mp \operatorname{csch} A^{-1} k_0 z, \quad (3.11)$$

the sign associations for fixed A depending upon the angular range in which β falls. From (3.11), also,

$$\cos\beta = -\tanh A^{-1}k_0Z. \quad (3.12)$$

Hence, from eqs. (3.7), (3.9), (3.11) and (3.12),

$$g''(\beta) = A \sinh A^{-1}k_0Z \cdot \cosh A^{-1}k_0Z, \quad (3.13)$$

which, from inspection of eq. (3.4), suggests either $A = 2$ or $A = -2$. Each choice gives

$$g''(\beta) = +\sinh k_0Z, \quad |A|=2, \quad (3.14)$$

and eq. (3.4) becomes

$$\begin{aligned} \ddot{z} - (2C - \dot{x}^2) k_0 \sinh k_0Z \\ = \mp 2a \sinh k_0x \cdot \sinh \frac{1}{2} k_0Z; \end{aligned} \quad (3.15)$$

from eq. (2.13)

$$\begin{aligned} \ddot{x} + \frac{1}{2} b \operatorname{sech}^2 \frac{1}{2} k_0Z \cdot \sinh 2k_0x \\ = \mp a \operatorname{sech} \frac{1}{2} k_0Z \cdot \cosh k_0x. \end{aligned} \quad (3.16)$$

The "energy equation", eq. (2.17), reads

$$\begin{aligned} \dot{x}^2 + \dot{z}^2 \pm 2ak_0^{-1} \sinh k_0x \cdot \operatorname{sech} \frac{1}{2} k_0Z \\ + bk_0^{-1} \sinh^2 k_0x \cdot \operatorname{sech}^2 \frac{1}{2} k_0Z = 2C, \end{aligned} \quad (3.17)$$

where

$$\dot{z}^2 = \frac{1}{4} \operatorname{sech}^2 \frac{1}{2} k_0Z \cdot \dot{Z}^2. \quad (3.18)$$

The upper and lower signs in eqs. (3.15) - (3.17) correspond to those in eqs. (3.11). Notice, finally, the periodic functions of z now are gone.

There is a consideration of domain, and we have a choice of procedures in defining the transformation further. Eqs. (3.2) and (3.9) give

$$\dot{z} = \epsilon(A) \cdot \frac{1}{2} \sin k_0 z \cdot \dot{Z} \quad (3.19)$$

while also, by definition of g ,

$$\frac{dZ}{dz} = g' = \epsilon(A) / \frac{1}{2} \sin k_0 z. \quad (3.20)$$

The second equation shows that Z , which takes all values as β varies between any two successive integer multiples of π (cf. eq. (3.10)), rises or falls monotonically with z , for fixed A , according to the sign of $\sin k_0 a$. We may choose $A > 0$ for $\sin \beta > 0$ and $A < 0$ for $\sin \beta < 0$. In this case also, eq. (3.19) shows that \dot{z} and \dot{Z} have the same sign always. There seems to be no special advantage to doing this, however, so I will just take $A = +2$ and allow Z to reverse sign relative to \dot{z} each time β passes through an integer multiple of π . Each time it does the upper (lower) signs in eqs. (3.11), and (3.15) - (3.17) give way to the lower (upper) signs. For finite \dot{z} both Z and \dot{Z} are infinite across these points, while x and \dot{x} are continuous. The system of equations of the axial motion along the z -axis from $-\infty$ to $+\infty$, coupled with the x -motion, has been replaced by a pair of systems corresponding to motion along succeeding half-cycles of $k_0 z$, covering a combined finite length $2\pi/k_0$. And in Z -space, there are axially reflecting mirrors at $\pm \infty$. The situation is illustrated in Figure 2.

For vanishing P_y , $a = 0$ and eqs. (3.15) and (3.16) become slightly simpler,

$$\ddot{Z} - (2C - \dot{x}^2) k_0 \sinh k_0 Z = 0 \quad (3.21)$$

$$\ddot{x} + \frac{1}{2} b \operatorname{sech}^2 \frac{1}{2} k_0 Z \cdot \sinh 2k_0 x = 0. \quad (3.22)$$

Note, once more, the absence of the periodic functions of z .

CHAPTER IV

CONCLUSION

A planar undulator wiggler field configuration has been constructed and examined from the standpoint of a transformation of the axial coordinate that sends the description of the single-electron motion between $z = \pm \infty$, to that of a particle in a box in which an oscillation cycle between $Z = \pm \infty$ corresponds to the traversal of a single cycle of wiggler period $|\Delta z| = 2\pi/k_0$.

A similar transformation is possible for the x-motion, but this has not been discussed as the resulting system does not appear especially interesting.

The special case $P_y = 0$ should be an interesting one physically (cf. remark in the paragraph following eq. (2.12)), so eqs. (3.21) and (3.22), in particular, might repay some effort of further investigation.

ACKNOWLEDGMENTS

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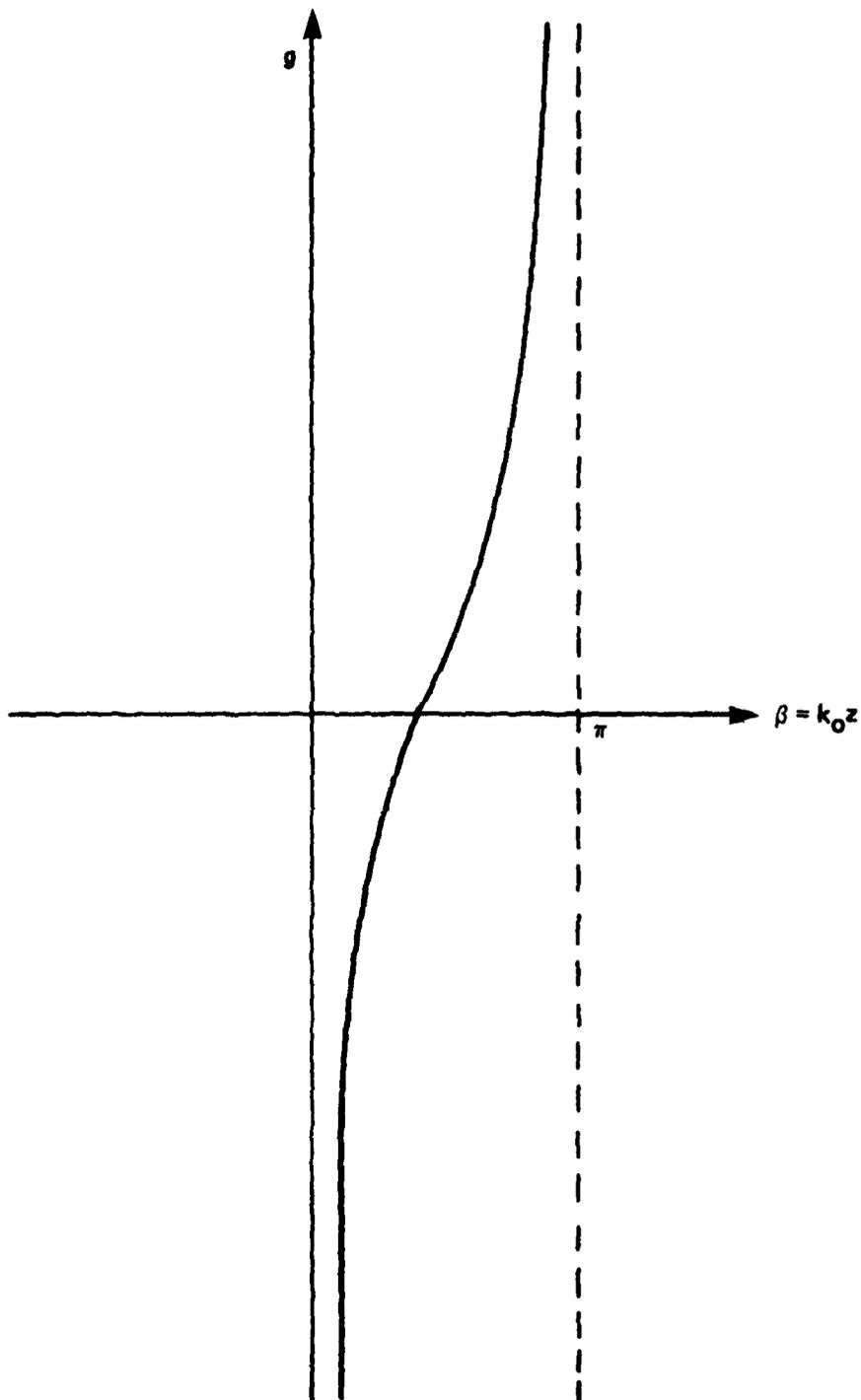


FIGURE 1 PRINCIPAL BRANCH OF $g(\beta)$ FOR $A > 0$ AND $B = 1$

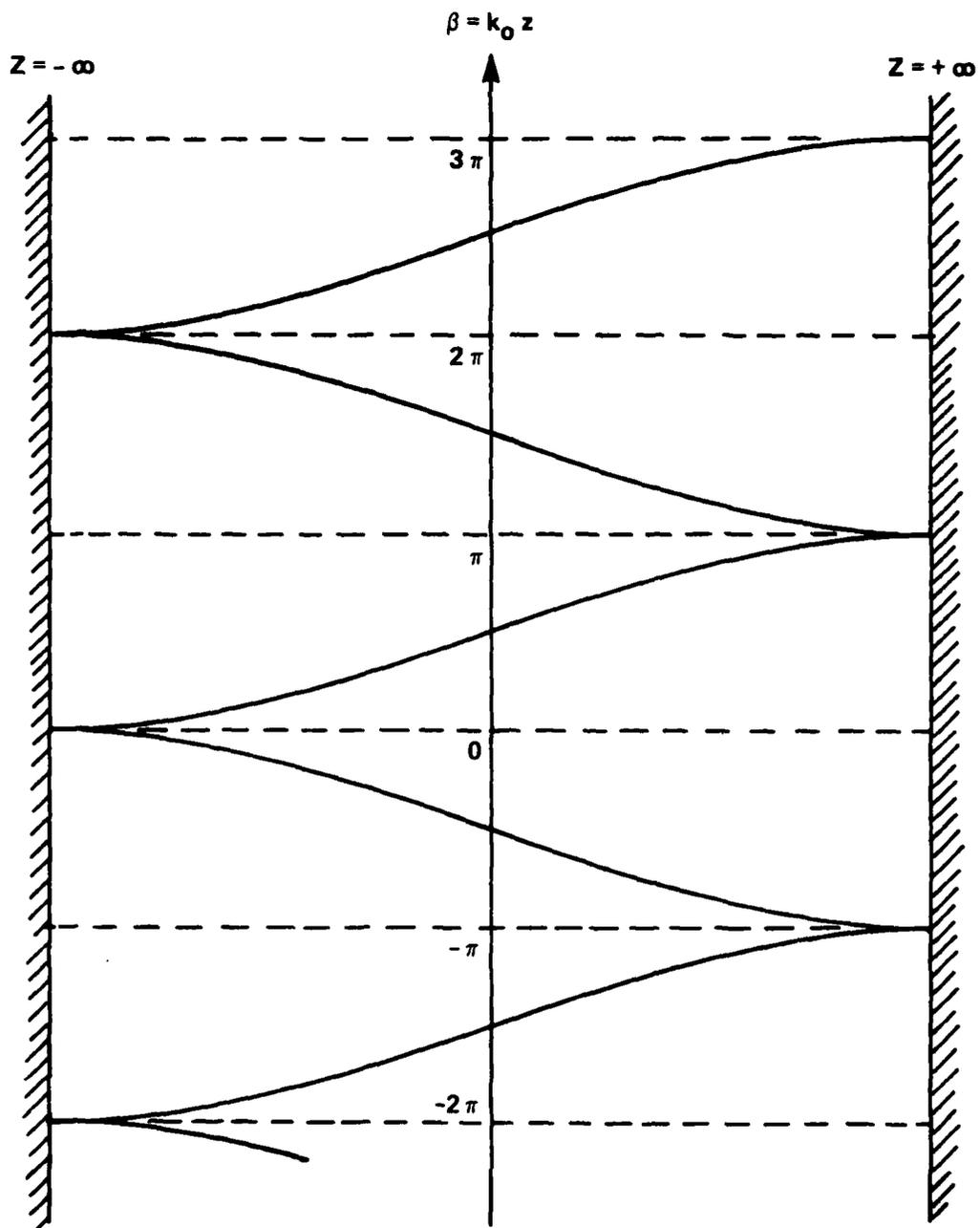


FIGURE 2 SKETCH OF $k_0 z$ VS Z FOR $A = +2$

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