BREAKDOWN OF LINDSTEDT EXPANSION FOR CHAOTIC MAPS

GUIDO GENTILE AND TITUS S. VAN ERP

Abstract. In a previous paper of one of us [Europhys. Lett. 59 (2002), 330–336] the validity of Greene's method for determining the critical constant of the standard map (SM) was questioned on the basis of some numerical findings. Here we come back to that analysis and we provide an interpretation of the numerical results, by showing that the conclusions of that paper were wrong as they relied on a plausible but untrue assumption. Hence no contradiction exists with respect to Greene's method. We show that the previous results, based on the expansion in Lindstedt series, do correspond to the critical constant but for a different map: the semi-standard map (SSM). For such a map no Greene's method analogue is at disposal, so that methods based on Lindstedt series are essentially the only possible ones. Moreover, we study the expansion for two simplified models obtained from the SM and SSM by suppressing the small divisors. We call them the simplified SM and simplified SSM, respectively; the first case turns out to be related to Kepler's equation after a proper transformation of variables. In both cases we give an analytical solution for the radius of convergence, that represents the singularity in the complex plane closest to the origin. Also here, the radius of convergence of the simplified SM turns out to be lower than that of the simplified SSM. However, despite the absence of small divisors these two radii are lower than those of the true maps (i.e. of the maps with small divisors) when the winding number equals the golden mean. Finally, we study the analyticity domain and, in particular, the critical constant for the two maps without small divisors. The analyticity domain turns out to be a perfect circle for the simplified SSM (as for the SSM itself), while it is stretched along the real axis for the simplified SM, yielding a critical constant which is larger than its radius of convergence.

1. Introduction

The Taylor-Chirikov map [12, 23] or standard map (SM) is one of the best known nonlinear models showing the onset of chaos in Hamiltonian systems. It describes with some level of approximation many physical systems. Among these there are numerous applications to plasma-physics, the field in which it was originally introduced. The SM is also exactly related to the time evolution of the "kicked rotor" and the equilibrium condition for a chain of masses superpositioned on a periodic potential. The latter model is known as the Frenkel-Kontorova (FK) model. This model is of equal importance for solid state physics as the SM is for plasma physics. It has, e.g., been applied to Josephson junctions arrays, charge density waves and surface friction [19]. More importantly, due to their simplicity and, yet, the complex behaviour they show, these minimalistic models have had an enormous impact for the understanding in complex phenomena such as nonlinearity, chaos, quasi-periodicity, and commensurate-incommensurate transitions. Although now part of any text-book in nonlinear physics and studied extensively over many years, the SM and FK still bear many unsolved problems. The most intriguing one of these is the sudden transition from smooth to chaotic orbits in the SM when the coupling parameter K is increased above a critical value K_c. In the FK model this transition is connected to change from a sliding to a pinned state and bears the name of Aubry transition (or analyticity breaking transition).

Key words and phrases. Standard map; semi-standard map; Frenkel-Kontorova; perturbation theory; critical constant; natural boundary.

The theoretical framework that characterizes this transition originates from the Kolmogorov-Arnol′d-Moser (KAM) theorem [2], that deals with the problem of small divisors that can occur in any perturbation expansion for quasi-integrable systems. In fact, the KAM theorem can be used to prove the non-chaotic behaviour of the SM for very small coupling K and sufficiently irrational winding number ω. Other arguments can then be applied to prove that a chaotic regime exists for values of K large enough giving an upper bound to K_c. For ω equal to the golden mean there exists an analytical bound by Mather, $K_c < 4/3$ [34], and the computer assisted proof of MacKay and Percival $K_c < 63/64 \approx 0.9844$ [33]. Moreover, another computer assisted analysis of Jungreis excluded the value $K = 0.9718$ for possible occurrence of invariant circles (smooth orbits) [27].

There exist several methods to calculate K_c precisely, among which Greene's method [23] has shown to be one of the most effective giving the estimate $K_c = 0.971635$. This method is based on the assumption that the dissolution of invariant curves can be associated with the sudden change from stability to instability of nearby closed orbits. The renormalization technique of MacKay (cf. Ref. [32], §4.4.1) is a further refinement of this method and has established the same value (cf. p. 199 of the quoted reference) with higher digit precision with respect to the original Greene's result. Yet, Greene's hypothesis has only been partly proven. A result by Falcolini and de la Llave [16] and, independently, by MacKay [31] yields that the critical constants for symplectic maps can never be higher than the ones obtained by Greene's method. Recently, the result has been extended to nontwist maps by Delshams and de la Llave [14]. Hence, $K_c \leq 0.971635$ for the SM with the golden mean as winding number. We mention that MacKay also showed that Greene's method does not apply to every map [31]. However, in the case of the SM, the situations for which no rigorous result can be given are considered unlikely; cf. the discussion in Refs. [16, 31].

Another way to calculate (or, at least, to estimate) K_c is through the Lindstedt series expansion. Any smooth invariant curve in the SM can be described, for complex K small enough, say $|K| < \rho$, where ρ will depend on the winding number of the curve, by an analytic function which conjugates the dynamics to the unperturbed one. Of course ρ provides a lower bound for K_c, which is essentially the maximum real value of K for which there is an analytic invariant curve with the fixed winding number (more precise definitions will be given below). By writing down the Taylor expansion and equating the Taylor orders in the functional equation satisfied by this conjugation function, the Fourier-Taylor coefficients can, in principle, be derived from the ones of lower order. In Ref. [42] an evaluation of this expansion, always in the case of the golden mean, to high orders led the authors to infer a convergence to a value $K_c \sim 0.97978$, which is higher than Greene's result[1]. In this article, we revisit that analysis and show that an apparent plausible assumption made in Ref. [42] is falsified beyond Taylor order $n > 200$. As a result, the Lindstedt expansion does not contradict Greene's result. The value $K_c \sim 0.97978$, however, does correspond to the critical value for a different map, the semi-standard map (SSM); cf. also Ref. [24]. We come back to this in Sect. 3; though, we not since now that this method is fundamental for the SSM, where no analogue of Greene's method exists (see also the comments at the end of Sect. 2).

Aubry [3] proposed another method, which is probably not very effective for high precision evaluation in a computer algorithm, but still interesting. It is based on an eigenvalue calculation of the dynamical matrix for the FK chain close to the critical point. Although this, in principle, requires the diagonalization of an infinite matrix, one can use the fact that the eigenvector of the lowest mode tends to localize [43]. The instability of the FK chain can than be determined in successive approximants by calculating the determinants of finite matrices of increasing size.

[1]Note that the fact that ρ and K_c are different can not invoked to explain the discrepancy, because of the direction of the inequality between the two quantities. One could also wonder what information about K_c can be inferred from ρ other than a lower bound. In general none, but one has strong numerical evidence that $\rho = K_c$ for the golden mean.

Another effective method is the frequency analysis method proposed by Laskar, which has been applied to the standard map in Ref. [30], giving for the golden mean a critical value $K_c = 0.9718$, which is close to Greene's value. In Ref. [11] it has been used to study numerically the dependence of the critical constant on the winding number.

Finally we mention the use of Padé approximants[2] to study numerically the entire analyticity domain. This is a powerful numerical tool even if is it less precise than other methods for detecting the critical constant K_c and not completely under control from a rigorous point of view. It has, for instance, been employed in Ref. [4] and, very recently, in Ref. [5], where the existence of a natural boundary for the analyticity domain of the SM has been checked numerically. Always with the aim of studying the analyticity domain Falcolini and de la Llave [17] developed a variant of Greene's method working for complex values of the parameter K that gives an alternative to the Padé approximants approach. An implementation of Padé approximants will be given below (in Sect. 4), though for a case in which the analytical solution is known.

Eventually, these approaches are assumed to converge to the same value. However, the proof of this is highly non-trivial. The ultimate goal, of course, would be to gain an analytical expression for K_c. This is still far beyond our capabilities. Inspired by the desire to investigate further the influence of the small divisors in the Lindstedt series expansion, we introduce two simplified models by setting rigorously all the divisors equal to 1 both for the SSM and the SM. In the latter case, this is a very well known model, Kepler's equation [50], which turns out to have a very similar transition and can be solved analytically. The radii of convergence are found to be lower than those determined with the methods described above for the SM and SSM, respectively, in case of golden mean winding numbers.

This article is organized as follows. In Sect. 2 we recall the definition of the SM and SSM. In Sect. 3 we come back to the analysis of Ref. [42] showing that, contrary to what was asserted in that paper, the Lindstedt expansion does not violate Greene's method, and we make the comparison between the SM and SSM. In our opinion the analysis in Sect. 3 gives some insight onto the mechanism of break up of the invariant curves, and leaves some open problems: this will be discussed to more extent in Sect. 5. In Sect. 4 we present a new model in which we suppress the small divisors and give an analytical expression both for the radius of convergence and the critical constant. Moreover, the analysis in Sect. 4 has some consequence on the cases with small divisors in relation with the appearance of a natural boundary in the analyticity domain; this is further discussed in Sect. 5. Finally we end up with the conclusions in Sect. 5.

2. The (Semi-) Standard Map

The SM and SSM can be written as

$$\begin{pmatrix} x_{i+1} \\ x_i \end{pmatrix} = \mathrm{T} \begin{pmatrix} x_i \\ x_{i-1} \end{pmatrix} = \begin{pmatrix} 2x_i + V'(x_i) - x_{i-1} \\ x_i \end{pmatrix}, \tag{1}$$

with $\{x_i\}$ defined mod 1, and

$$V'(x) = \begin{cases} \dfrac{K}{2\pi} \sin(2\pi x) & \text{for the SM,} \\ \dfrac{K}{4\pi i} \exp(i 2\pi x) & \text{for the SSM.} \end{cases} \tag{2}$$

The resulting sequence $\{x_i \bmod 1\}$, for $i = 2, \ldots \infty$, originating from a starting point (x_1, x_0) corresponds to a discrete trajectory on an invariant curve, when the latter exists. Such a trajectory for the SM can be related to the equilibrium positions of an infinite FK chain

[2]The (L, M)-Padé approximant for a function $f(x)$ is given by the ratio of two polynomials, $f(x) \approx P_L(x)/Q_M(x)$, with $P_L = p_0 + p_1 x + \ldots + p_L x^L$ and $Q_M = 1 + q_1 x + \ldots q_M x^M$. Hence, the simple Taylor expansion of order n can be considered as a special case of Padé approximant with $L = n$ and $M = 0$.

where particles with harmonic nearest neighbour coupling are placed on a periodic potential $V(x) = K(2\pi)^{-2}(1 - \cos(2\pi x))$. The SSM has not a similar counterpart, but is much simpler in its mathematics, and this is basically the reason why it was considered by Greene and Percival [24]. By definition, in terms of the lift of the map, the *winding number* or *rotation number* of an orbit is given by

$$\omega \quad \equiv \quad \langle x_{i+1} - x_i \rangle \equiv \lim_{n\to\infty} \sum_{i=0}^{n} \frac{x_{i+1} - x_i}{n} = \lim_{n\to\infty} \frac{x_n - x_0}{n}, \tag{3}$$

when the limit exists. For low coupling K and ω incommensurate to the periodicity of V', there exists a continuous function $g(x; K, \omega)$ such that the positions $\{x_i \bmod 1\}$ can be expressed as $x_i = g(i\omega+\varphi; K, \omega)$, where φ is an arbitrary phase. This function is often called the *conjugating function* or, in context with the FK model, the modulation or hull function. Its shape depends on the winding number ω and on the coupling parameter K.

The conjugating function satisfies the functional equation

$$2g(x; K, \omega) - g(x + \omega; K, \omega) - g(x - \omega; K, \omega) = -V'\big(x + g(x; K, \omega)\big). \tag{4}$$

For K large enough the function $g(x; K, \omega)$ becomes discontinuous. For the SM this implies that the orbits become chaotic and for the FK that the chain of particles gets pinned together with the appearance of a phonon gap.

There are several quantities of interest which one can introduce in order to study the transition from regular to chaotic dynamics. As the function $g(x; K, \omega)$ is analytic for K close to the origin one can consider its series expansion in powers of K,

$$g(x; K, \omega) \quad = \quad \sum_{n=1}^{\infty} K^n g^{(n)}(x; \omega), \tag{5}$$

and define the *radius of convergence* $\rho(\omega)$ as

$$\rho(\omega) \quad = \quad \inf_{x\in[0,1]} \Big(\limsup_{n\to\infty} \Big| g^{(n)}(x; \omega) \Big|^{1/n} \Big)^{-1}. \tag{6}$$

Note that the infimum appears in the definition of the radius of convergence because, as a result of the incommensurate winding number ω, each invariant curve is filled densely by any trajectory lying on it. Hence, existence of the invariant curve itself requires the latter to be defined for all $x \in [0, 1]$.

The *critical constant* is defined as the (positive) real value $K_c(\omega)$ such that for $K > K_c(\omega)$ the conjugating function is not analytic any more[3]. It is believed that the analyticity domain of the conjugating function has a natural boundary [24, 4, 5]: this means that $g(x; K, \omega)$ has a set of singularities in terms of K that form a closed curve around the origin in the complex plane. Hence, the radius of convergence $\rho(\omega)$ corresponds to the singularity closest to the origin, while the critical constant $K_c(\omega)$ corresponds to the intersection of this curve with the (positive) real axis. By definition one has $K_c(\omega) \geq \rho(\omega)$, so that by estimating the radius of convergence one finds a lower bound for the critical constant. Furthermore, it is generally accepted that $K_c(\tau) = \rho(\tau)$ for the *golden mean*[4] $\tau = (\sqrt{5} - 1)/2 \approx 0.618034$, whereas there is strong numerical evidence that $K_c(\omega)$ can be much larger than $\rho(\omega)$ for winding numbers which have very large partial quotients in their continued fraction expansion[5] [11, 9]. It is also

[3]The reason why one usually does not considers the negative critical constant, that is the negative value $K_c'(\omega)$ such that for $K < K_c'(\omega)$ there is no longer an analytic invariant curve, is that $K_c'(\omega) = -K_c(\omega)$ for the SM.

[4]The golden mean is sometimes in other literature defined as the inverse of this value: $(\sqrt{5}+1)/2 = \tau^{-1} \approx 1.618034$.

[5]For an introduction of the continued fraction theory and a discussion of the basic properties we refer to the classical textbook by Hardy and Wright [26]. In the continued fraction expansion of a number $\omega =$

commonly believed (on the basis of numerical simulations and heuristic arguments) that $K_c(\omega)$ has the highest value for the golden mean $\omega = \tau$.

So far, the most accurate method to calculate K_c is based on Greene's method (also known as residue criterion). In this method the infinite trajectory $\{x_i \bmod 1\}$ with irrational winding number ω is approached by successive approximants which are periodic trajectories with rational winding numbers $\omega_j = p_j/q_j$ tending to ω, and $x_{i+q_j} \bmod 1 = x_i$. Hence, p_j and q_j are at each level j two integer values whose ratio gives a better estimate of ω for each increment in j and $\omega = \lim_{j\to\infty} \omega_j$. These numbers can, for instance, be obtained using the best approximants in the continued fraction expansion of ω. For $\omega = \tau$ this results in the ratios of subsequent Fibonacci numbers ($\tau \approx F_{j-1}/F_j$ with $F_0 = F_1 = 1$ and $F_j = F_{j-1} + F_{j-2}$ for $j > 1$). Conclusively, Greene's method tells how to construct the periodic orbits and to measure their stability by means of a suitable quantity, called the *residue*, which does not tend to zero any more for $K > K_c$.

Besides being only partly proven, Greene's method has also some other limitations. For instance, this method does not work for other interesting models, as the SSM and Siegel's problem [39], where the construction of periodic orbits fails. One can also easily check the non-existence of smooth periodic orbit by a first orders perturbation theory. The best general alternative is the Lindstedt series expansion. This method is more generally applicable (it also works for the SSM and for any Hamiltonian systems close to an integrable one), but, in view of a numerical implementation, is less accurate than Greene's method for the SM and, in general, is more suitable for studying the radius of convergence rather than the critical constant.

3. The Lindstedt series expansion

3.1. **Standard Map.** A way to study the transition is by means of the *Lindstedt series*, which in this case means the expansion of the function $g(x; K, \omega)$ both in Fourier and in Taylor series. Such expansions were originally introduced by Lindstedt and Newcomb to study problems in celestial mechanics [37]. By defining the Fourier transform as

$$\begin{aligned} g(x; K, \omega) &= \sum_{k=-\infty}^{+\infty} \hat{g}_k(K, \omega)\, \mathrm{e}^{2\pi i k x} \quad \text{with inverse:} \\ \hat{g}_k(K, \omega) &= \int_0^1 \mathrm{d}x\, g(x; K, \omega)\, \mathrm{e}^{-2\pi i k x}, \end{aligned} \tag{7}$$

and expanding

$$\hat{g}_k(K, \omega) = K \hat{g}_k^{(1)}(\omega) + K^2 \hat{g}_k^{(2)}(\omega) + K^3 \hat{g}_k^{(3)}(\omega) + \ldots, \tag{8}$$

we end up with Fourier-Taylor coefficients $\hat{g}_k^{(n)}(\omega)$, where n is the Taylor index and k is the Fourier index. Of course, $\hat{g}_k^{(n)}(\omega)$ depends on ω, but henceforth we withdraw such a dependence in order not to overwhelm the notation, whenever no ambiguity can arise.

Now, using Eq. (4) we can relate the Fourier-Taylor coefficients of order n by the ones with lower Taylor index by [42]

$$\begin{aligned} D_k^2\, \hat{g}_k^{(n)} = \frac{i}{4\pi}\Big\{\delta_{1,k} - \delta_{-1,k}\Big\}\delta_{1,n} + \frac{i}{4\pi} \sum_{m=1}^{\infty} \frac{(i2\pi)^m}{m!} \sum_{n_1+n_2+\ldots+n_m=n-1} \\ \Big\{ \sum_{k_1+k_2+\ldots+k_m=k-1} \hat{g}_{k_1}^{(n_1)} \hat{g}_{k_2}^{(n_2)} \ldots \hat{g}_{k_m}^{(n_m)} - (-1)^m \sum_{k_1+k_2+\ldots+k_m=k+1} \hat{g}_{k_1}^{(n_1)} \hat{g}_{k_2}^{(n_2)} \ldots \hat{g}_{k_m}^{(n_m)} \Big\}, \end{aligned} \tag{9}$$

$[a_0, a_1, a_2, \ldots] = a_0 + 1/(a_1 + 1/(a_2 + 1/(\ldots)))$, the numbers a_j are called the partial quotients, while the rational numbers $q_j/p_j = [a_0, a_1, a_2, \ldots, a_j]$ are called the best approximants for ω.

with

$$D_k^2 \equiv \frac{1}{\hat g_k(K,\omega)}\int_0^1 \mathrm{d}x\,\big(2g(x;K,\omega)-g(x+\omega;K,\omega)-g(x-\omega;K,\omega)\big)\,\mathrm{e}^{-2\pi ikx} \quad (10)$$
$$= 2\big(1-\cos(2\pi k\omega)\big) = \big(2\sin(\pi k\omega)\big)^2,$$

and where $\sum_{n_1+n_2+\ldots+n_m=n\pm1}$ implies a summation of all possible integers $n_1, n_2, \ldots, n_m$ with the constraint that $\sum_{i=1}^m n_i = n \pm 1$. There are ways to reduce the number of summations in Eq. (9). One possible way was proposed in Ref. [42] to construct an extended matrix $P(n,k,m)$ defined as

$$P(n,k,m) = \frac{(2\pi i)^m}{m!}\sum_{n_1+n_2+\ldots+n_m=n}\ \sum_{k_1+k_2+\ldots+k_m=k} \hat g_{k_1}^{(n_1)}\hat g_{k_2}^{(n_2)}\ldots\hat g_{k_m}^{(n_m)}. \quad (11)$$

One can show that $P(n,k,m)=0$ if $|k|>n$ or $m>n$. This gives rise to the following recursive relations [42]:

$$P(1,\pm1,1) = \frac{\mp1}{2D_1^2}, \quad (12)$$
$$P(n,k,1) = -\frac{1}{2}D_k^{-2}\sum_{m=1}^{n-1}\Big[P(n-1,k-1,m)-(-1)^m P(n-1,k+1,m)\Big],$$
$$P(n,k,m) = \frac{1}{m}\sum_{n'=1}^{n-m+1}\ \sum_{k'=\max\{-n',k-n+n'\}}^{\min\{n',k+n-n'\}} P(n',k',1)P(n-n',k-k',m-1), \quad 1<m\le n,$$

from which we can distract the Fourier-Taylor coefficients by

$$\hat g_k^{(n)} = \frac{P(n,k,1)}{2\pi i}. \quad (13)$$

The entries of P are all real and obey the symmetry relation $P(n,k,m)=(-1)^m P(n,-k,m)$. Moreover, besides being zero for $|k|>n$ and $m>n$, $P(n,k,m)$ has zero values whenever $k+n$ is odd. Hence, $k=n,n-2,\ldots,-n$ are the only non-zero entries of P.

The relations of Eqs. (12) are very efficient to evaluate $\hat g_k^{(n)}$, and they were used in Ref. [42] to reach a Taylor order of approximately $n=200$. To go beyond this limit, sufficient computer power and time is needed as both the computation time and the number of non-zero matrix entries increase with $\sim n^3$. Hence, memory can become a severe problem as the number of entries that have to be stored can easily go beyond the maximum allowed allocation limit of the computing system. Also the precision has to be high enough in order to minimize numerical errors, but this problem is easily solved without requiring a precision as high as in the case of winding numbers close to rational numbers, as in Ref. [5], where the small divisors could become really small and up to 480 digits were needed. In this work, we reached the level $n=700$ (see Fig. 1) and we believe that going beyond this order is not very profitable for obtaining a more accurate evaluation of K_c. We come back to these results after addressing the small divisor problem that arises from Eq. (12).

From Eq. (10) and the second line in Eq. (12) one sees that even for irrational values of ω, the terms D_k^{-2} can become arbitrarily high for some k. This effect is a typical example of the *small divisor problem* (or small denominator problem), that can strongly prevent the convergence of any perturbative series. In fact, in general it requires a stronger condition than irrationality, such as a *Diophantine condition*[6] [2]. Among all the irrational numbers, the golden mean τ suffers the least from the small divisor problem and has therefore the highest convergence radius

[6]The usual Diophantine condition requires $|\omega q-p|>1/C_0|q|^\gamma$ for all $(p,q)\in\mathbb{Z}^2$ with $q\neq0$ and for suitable positive constants C_0 and γ. But one can require also the Bryuno condition, which is a condition stronger than irrationality but weaker than the usual Diophantine condition; cf. for instance Refs. [8, 21] in the case of the SM.

$K_c(\tau)$. The golden mean is "relatively difficult" to approximate by rational numbers as those arising, for instance, from the continued fraction expansion[7] (one can say that it is the "most irrational number"). As addressed above, the best approximants for the golden mean are given by the ratios F_j/F_{j+1}, where $\{F_j\}$ is the sequence of the Fibonacci numbers. Therefore, the small divisors will be largest for $k = F_j$. From the exact relation $F_{j-1} - F_j\tau = (-\tau)^{j+1}$, one can show that for $j \to \infty$

$$D_{F_j}^{-2} \approx \frac{1}{4\pi^2}\Big(\frac{1}{\tau^2}\Big)^{j+1} \sim (2.618)^{j+1}. \tag{14}$$

However, we would like to stress that the small divisor problem is not the only mechanism causing the breakdown of the perturbative approach. This becomes evident in Sect. 4 where we introduce the model that arises when we rigorously set $D_k^{-2} = 1$ for all k in the series of Eq. (12). Clearly, this simplified expansion can not be affected by the small divisors. However, it still has a radius of convergence and a critical constant, as shown by the analytical solution. As the radius of convergence ρ of this simplified model is found to be lower than $\rho(\tau) = K_c(\tau)$ for the SM, it proves that the golden mean winding number is remarkably resistant to the problem of small divisors. The full analysis of this model is given in Sect. 4.

Coming back to the results of Fig. 1, we see that indeed the evolution of $P(n, k, 1)$ makes sudden jumps at the Fibonacci numbers as expected from Eq. (14). Besides, deceptively the values $n = k = F_j$ seem to determine the whole power law behaviour of $\hat{g}_k^{(n)}$, which is true till Taylor order $n \sim 200$. This assumption made in Ref. [42] allows for a further simplification of Eq. (12) by defining the reduced matrix $Q(n, m) \equiv P(n, n, m)$ obeying the relations

$$\begin{aligned} Q(1,1) &= \frac{-1}{2D_1^2}, \\ Q(n,1) &= -\frac{1}{2}D_n^{-2}\sum_{m=1}^{n-1} Q(n-1,m), \\ Q(n,m) &= \frac{1}{m}\sum_{n'=1}^{n-m+1} Q(n',1)Q(n-n',m-1), \qquad 1 < m \le n. \end{aligned} \tag{15}$$

This set of equations make high order ($\sim n = F_{20} = 10946$) evaluations accessible for computer calculations. At this order the value of the radius of convergence seems to stabilize near $\rho = 0.97978$, but this is still higher than the one obtained by Greene's method. As a consequence, validity of Greene's method was questioned in the quoted paper [42].

The more elaborated calculations in this work show that the assumption made in Ref. [42] was actually wrong as shown by the high order behaviour in Fig. 1. Still, we find that $k = n$ gives the maximum for $|P(n, k, 1)|$ whenever n is a Fibonacci number. However, the character of the evolution changes from being peaked to more smooth oscillations. Clearly, the line connecting the points $(F_j, |P(F_j, F_j, 1)|)$ does no longer dominate the increment of the entries of P for $n > 200$. As $|P(n, k_{\max}(n), 1)|^{1/n}$ shows local maxima for $(n, k_{\max}) = (383, 377)$ and $(n, k_{\max}) = (622, 610)$ just after F_{13} and F_{14}, we fitted the line $\alpha_2\lambda_2^n$ through the corresponding points. From this fit, $\lambda_2 = 1.0248$, the estimate for $K_c \approx \lambda_2^{-1} = 0.9758$ is obtained. Although still higher than Greene's value, it is already considerably lower than $\rho = 0.97978$ obtained from Eqs. (15) for $n = F_{20} = 10946$. Note that the latter approach of Ref. [42] for this lower approximant $n = F_{14} = 610$, as obtained from the line λ_1^n (See Fig. 1), would result in $\rho \approx 1/\lambda_1 = 0.9817$, still approximately 0.002 higher than the nearly converged value of $\rho = 0.97978$. Hence, a decay of 0.004 from 0.9758 at $n = 622$ to 0.9716 at $n \to \infty$ is not unlikely. As a consequence, contrary to the results of the restricted series (15), there is no evidence at all that the full Lindstedt series (12) violates Greene's hypothesis. This also shows

[7] A more formal statement is that τ allows the smallest values of C_0 and γ in the Diophantine condition (see previous footnote). Hence, it is the number that is the most distant from all rational numbers.

that a further simplification of Eqs. (12) is not easily obtained and that an accurate evaluation of K_c based on the Lindstedt perturbation is severely demanding.

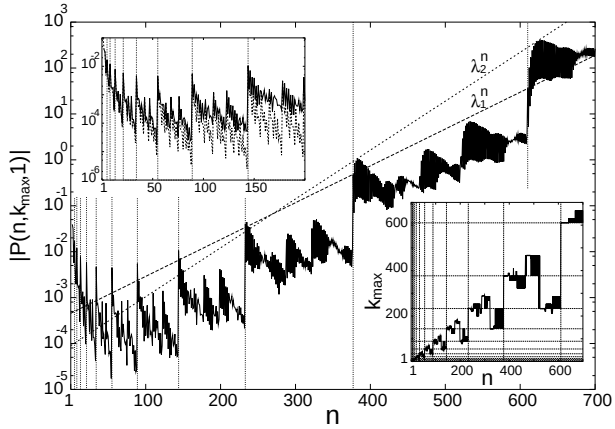

FIGURE 1. $|P(n, k_{\max}, 1)|$ as a function of n. This is defined as the maximum value of $|P(n, k, 1)|$ of all k; hence, $k_{\max} = k_{\max}(n)$ is defined as the k value where $|P(n, k, 1)|$ has this maximum. The inset in the lower corner shows $k_{\max}$ as function of n. The inset in the left upper corner is an enlargement of the first 200 terms together with $|P(n, n, 1)|$ (dashed line). From these figures, one can clearly detect sudden boosts in the function $|P(n, k_{\max}, 1)|$ where $k_{\max} = n$ at the Fibonacci values (dashed vertical lines). However, whereas for $n < 200$ the character is sharply peaked at these values, its behaviour changes for higher orders. Still $k_{\max} = n$ for n a Fibonacci number, but the intersecting line described by $\alpha_1 \lambda_1^n$ does no longer dominate the complete evolution of all the $|P(n, k, 1)|$ terms. $\lambda_1 = 1.0186$ is determined by the line through $(n, k) = (F_{13}, F_{13}) = (377, 377)$ and $(n, k) = (F_{14}, F_{14}) = (610, 610)$. $\lambda_2 = 1.0248$ is set by the line through $(n, k) = (383, 377)$ and $(n, k) = (622, 610)$ where $|P(n, k_{\max}(n), 1)|^{1/n}$ shows local maxima in n. The inversed values, $\lambda_1^{-1} \approx 0.9817$ and $\lambda_2^{-1} \approx 0.9758$ are assumed to converge for higher n to K_c for the SSM and SM, respectively.

3.2. **Semi-Standard Map.** When evaluating Eq. (1) for the SSM (2), the factors $\delta_{-1,k}$ and $-(-1)^m \sum \dots$ are not present in Eq. (9). It is then straightforward to show that the only non-zero entries of $P(n, k, m)$ in (11) are those with $n = k$. Hence, the assumption made in Ref. [42] that gave rise to Eqs. (15) does not corresponds to the critical constant of the SM, but still gives the correct value for SSM. A result by Davie [13] shows that, for maps including the ones we are considering, the radius of convergence (6) is equal to

$$\rho(\omega) \quad = \quad \Big(\limsup_{n\to\infty} \max_{|k|\le n} \Big| \hat{g}_k^{(n)}(\omega) \Big|^{1/n} \Big)^{-1}. \tag{16}$$

Therefore, the radius of convergence for the SM cannot be larger than the radius of convergence of the SSM, but of course it implies only a lower bound on the critical constant. Numerically by using Padé approximants in Ref. [5] it has been found that for certain values of the winding number ω, the radius of convergence of the SM is strictly smaller than the radius for the SSM. For the golden mean it is hard to improve upon simple power series using Padé, but for

other numbers closer to resonant values it is possible and the phenomenon becomes much more evident.

The fact that the radius of convergence for these numbers is lower for the SM than for the SSM implies, by (16), that dominant contributions arise from terms with Taylor orders n for which $|\hat{g}^{(n)}_{k_{\max}}| > |\hat{g}^{(n)}_n|$, where $k_{\max} = k_{\max}(n)$ is defined as the value of k maximizing $|P(n,k,1)|$ at fixed n. This is exactly what emerges from the numerics as noted above and shown in Fig. 1 for $n > 200$. Clearly, this is not the case for the SSM where one can limit to $k = n$. Hence, in Ref. [42] $K_c(\tau) = \rho(\tau)$ was actually determined for the SSM to be 0.97978 at Taylor order $n = F_{20} = 10946$. In the calculations of this work, we went to order $n = F_{24} = 75025$, that gave the value 0.97937. As the root criterion saturates very slowly, the numerical results provide essentially only an (accurate) upper bound for the radius of convergence.

To summarize, we found that, also for the golden mean, the radius of convergence for the SM is strictly less than the radius of convergence for the SSM. Therefore, as a general comment we can remark that for the SM the presence of all harmonics in the Fourier expansion of the Taylor coefficients, that is $g^{(n)}(x;\omega) = \sum_{k\in\mathbb{Z}} \hat{g}^{(n)}_k(\omega)\mathrm{e}^{2\pi i x}$, has a double effect. On the one hand the radius of convergence becomes smaller with respect to that of the SSM. On the other hand, the critical constant $K_c(\omega)$ can be larger than its radius of convergence $\rho(\omega)$. For the golden mean the two values are equal as emerges numerically [17], but for other values they can be appreciably different. One can imagine that the first phenomenon is due to the presence of contributions $\hat{g}^{(n)}_k(\omega)$ larger (in modulus) than $\hat{g}^{(n)}_n(\omega)$, while the second one is a consequence of deep cancellations between the harmonics of given perturbative order. These two effects are, in general, much more dominant for winding numbers ω close to rational values (see for instance Refs. [5, 6, 8, 9, 10]).

4. Putting small divisors to unity

4.1. **Introduction of the simplified maps.** An interesting study appears if we set rigorously all possible small divisors equal to unity: $D_k = 1$ for all k in Eqs. (12) and (15). Although the inspiration of this model was simply the study of the perturbation expansion when the small divisors have no effect, we can retrace from this series back to a functional relation as the one in Eq. (4) for a function $h(x;K)$

$$\begin{aligned} h(x;K) + \frac{K}{4\pi i}\exp\big(i2\pi\big(x + h(x;K)\big)\big) &= 0 \quad \text{for the SSM and} \\ h(x;K) + \frac{K}{4\pi}\sin\big(2\pi\big(x + h(x;K)\big) &= 0 \quad \text{for the SM,} \end{aligned} \tag{17}$$

where ω has vanished. Hence, the divergence of the simplified series corresponds to values K where the functional equations (17) have no analytical solution any more. A logical next step would be to relate Eqs. (17) to the iteration of a map similar to Eq. (1). As the relation (17) does no longer contain the arguments $x\pm\omega$ this is not so evident. However, one can relate the function $h(x;K)$ to the hull function of a FK-type system. It can be shown that this corresponds to a one-dimensional Einstein solid that is interacting with an external incommensurate potential[8]. Due to the lack of neighbour interaction, which makes each particle independent, it is highly unusual to describe for such a system the equilibrium coordinates by a collective hull function. Still, there are no restriction not to do so and one can even give such a function a physical meaning. As known from the FK model, the continuous shape of the hull function is directly associated with the existence of a sliding mode where the FK chain can slide over the periodic potential without cost of energy [19, 43]. In this case, the complete phonon spectrum is given by sum of oscillations of the individual particles that are not zero in general. The sliding mode appears when we add an extra degree of freedom to the system as shown in Fig. 2.

[8]The Einstein model is a well known approximation in solid physics where the vibrations of a lattice of N atoms is treated as a set of 3N independent harmonic oscillators in one dimension [29]

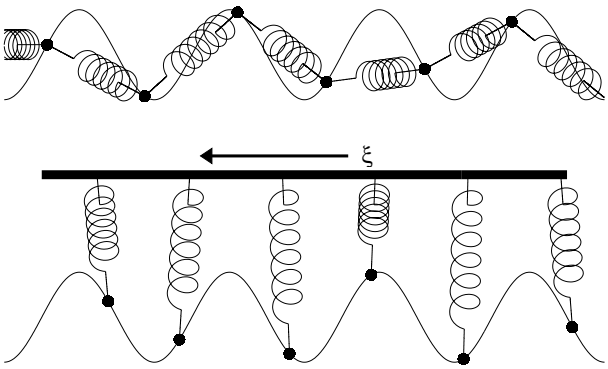

FIGURE 2. Illustration of the FK model (top) and the system that obeys Eq. (17) (bottom). The latter corresponds also to the FKT model without neighbour interaction. All particle are connected to the upper rod whose position is given by ξ. A sliding mode may exist when ξ can be varied.

Here, all particles have no interaction with their neighbours, but are connected to an upper rod. When the rod has an infinite mass compared to the particle masses, the system is basically an Einstein solid. However, if we assume that the position of the rod may vary according to a coordinate ξ an extra phonon mode exist that is zero for $K < K_c$ in this system. Hence, the breakdown of the Lindstedt expansion (12) with $D_k = 1$ for all k can be also related to a real physical sliding-pinning transition. The model as illustrated in Fig. 2 is also equal to a special case of the Frenkel-Kontorova-Tomlinson (FKT) model where each particle is connected to a rod by a spring with spring constant c_r and additionally to its neighbours with a coupling c_n[9]. The FKT model has been proposed to study more realisticly the frictional behaviour between atomic surfaces [45, 46, 25]. The model that is described by Eq. (17) simply corresponds to the FKT system with $c_r = 1$ and $c_n = 0$. The nice thing is that the perturbation series of Eqs. (17) can be solved exactly. To show this, we will start with the more simple SSM case.

4.2. **Radii of convergence for simplified maps.**

4.2.1. *Simplified SSM.* It can be convenient to use following normalization

$$R(n,m) \quad = \quad n!(-2)^n Q(n,m), \tag{18}$$

with matrix entries that are integer and positive and with $R(n,n) = 1$. From Eq. (12) and (18) with $D_n^{-2} = 1$ we derive

$$\begin{aligned} R(n,1) &= n\sum_{m=1}^{n-1} R(n-1,m), \\ R(n,m) &= \frac{1}{m}\sum_{n'=1}^{n-m+1}\binom{n}{n'}R(n',1)R(n-n',m-1), \qquad 1 < m \le n. \end{aligned} \tag{19}$$

[9]The simple Tomlinson model [41] is not the system in Fig. 2 as this usually consists of only one single oscillator (particle). This ancient model is now often applied to simulate the "stick-slip" motion of an atomic force microscope (AFM)-tip over a sample surface.

Note that the recursive relations in Eq. (19) for $m > 1$ coincide with those satisfied by the Stirling numbers of first and second kind, $S_n^{(m)}$ and $\mathfrak{S}_n^{(m)}$ respectively[10]. Of course what is different is the relation for $m = 1$.

From relations (19) with $R(1,1) = 1$ the following exact equality can be proven

$$R(n,m) \quad = \quad n^{n-m}\binom{n-1}{m-1}. \tag{20}$$

The proof of this equation is given in the appendices in two ways. In App. A.1 we derive this proof using the argument of induction. In App. B.1 we give a proof based on the tree formalism that was firstly introduced in Refs. [15, 36], then extended and formalized in Ref. [20]. The first proof is quite elementary, but a little elongated. The second is short, but less self-contained as it requires some knowledge of the previous publications about the tree formalism (but of course it becomes very simple for any reader acquainted with such a technique). Furthermore, the latter is most practical for the more complicated proof for the simplified SM in Eq. (17). Now, from Eq. (18) and (20) we deduce that

$$Q(n,1) \quad = \quad (-1)^n \frac{n^{n-1}}{2^n n!}, \tag{21}$$

which shows a power law behaviour $\sim a\lambda^n$ for large n giving the radius of convergence as $\rho = 1/\lambda$. Hence,

$$\ln|Q(n,1)| \quad \sim \quad (n-1)\ln(n) - n\ln(2) - \ln(n!). \tag{22}$$

Then by using[11] [47]

$$\ln(n!) \quad \approx \quad \left(n + \frac{1}{2}\right)\ln n - n + \frac{1}{2}\ln(2\pi), \tag{23}$$

we get

$$\begin{aligned} \ln|Q(n,1)| \quad &= \quad n(1-\ln(2)) - \frac{3}{2}\ln(n) - \frac{1}{2}\ln(2\pi) \\ \Rightarrow |Q(n,1)| \quad &\sim \quad \frac{1}{\sqrt{2\pi n^3}}\left(\frac{1}{2}\mathrm{e}\right)^n, \end{aligned} \tag{24}$$

yielding a radius of convergence $\rho = 2/\mathrm{e} \approx 0.735759$. This value is less than the SSM value or SM value. This can be a bit contra-intuitive, as one could have expected that the possible occurrence of small divisors would give a lower ρ. Apparently, this does not happen for the golden mean. This can be understood by the following reasoning. Although the small divisor factors D_k^{-2} can become arbitrary large for some k giving a boost to the series (12) and (15), at most values of k they will be considerable smaller than 1 resulting in an opposite effect. Hence, for the golden mean as winding number the latter effect seems to be more dominant yielding an even higher value for ρ than the case where all D_k^{-2} terms are equal to 1. We can say that, in fact, for the golden mean the small divisors "are not so bad", whereas they become really small for winding numbers much closer to rational values (that is with very large partial quotients in their continued fraction expansion).

[10] The Stirling numbers of the first kind $S_n^{(m)}$ are defined by the requirement that $(-1)^{n-m}S_m^{(n)}$ is the number of permutations of n symbols which have exactly m cycles. The Stirling number of the second kind $\mathfrak{S}_n^{(m)}$ is equal to the way of partioning a set of n elements into m non-empty subsets. See [1], p. 824, §24.1.3 and §24.1.4, with $r = m - 1$.

[11] This is a refinement of the well known Stirling's formula $\ln n! \approx n\ln n - n$.

4.2.2. *Simplified SM.* The simplified SM considered in Eq. (17) is well known in celestial mechanics [50] after applying the following variable transformation. Write $K/2 = -\mathfrak{e}$ and $2\pi x = \mathfrak{M}$ where $\mathfrak{e}$ is the eccentricity and $\mathfrak{M}$ is the mean anomaly. Then the eccentric anomaly $\mathrm{E} = 2\pi(x + h(x;K))$ is related to $\mathfrak{M}$ through Kepler's equation $\mathfrak{M} = \mathrm{E} - \mathfrak{e}\sin\mathrm{E}$, which is exactly the second equation in Eq. (17).

The recursive relations (12) with $D_k^{-2} = 1$ have also an exact solution that we write here:

$$P(n,k,1) \quad = \quad \begin{cases} \dfrac{(-1)^{n+(n-k)/2}}{2^n} \dfrac{k^{n-1}}{((n-k)/2)!((n+k)/2)!}, & \text{for } |k| \le n \text{ and } k+n \text{ is even} \\ 0, & \text{otherwise,} \end{cases} \tag{25}$$

which can be obtained by the Lagrange inversion theorem [48, 50]. We present a new derivation of this relation based on the tree formalism in App. B.2

Then, by using Eq. (16) we see that we have to compute the maximum over k of $|P(n,k,1)|$. By assuming that the maximum is reached for some k which is not too close to n[12] (an assumption that we shall verify a posteriori), we can approximate the factorial appearing in Eq. (25) with Stirling's formula (23). This gives rise to

$$\begin{aligned} |P(n,k,1)| \quad &\sim \quad \frac{1}{2^n} \frac{k^{n-1}\mathrm{e}^n}{(\frac{1}{2}(n-k))^{\frac{1}{2}(n-k)}(\frac{1}{2}(n+k))^{\frac{1}{2}(n+k)}} \frac{1}{(\frac{1}{4}(n^2-k^2))^{\frac{1}{2}}} \\ &\sim \quad \frac{2\mathrm{e}^n}{n^2} \frac{1}{\sigma(1-\sigma^2)^{\frac{1}{2}}} \left(\frac{\sigma}{(1-\sigma)^{\frac{1}{2}(1-\sigma)}(1+\sigma)^{\frac{1}{2}(1+\sigma)}} \right)^n, \end{aligned} \tag{26}$$

where we have defined $\sigma = k/n \in [-1,1]$. Hence, we have to compute the maximum of the function

$$E(\sigma) \quad = \quad \sigma \exp\left[-\frac{1-\sigma}{2}\ln(1-\sigma) - \frac{1+\sigma}{2}\ln(1+\sigma) \right]. \tag{27}$$

By taking the derivative $\partial E(\sigma)/\partial\sigma = 0$, we find that the maximum is reached at a value $\sigma_{\max}$ that satisfies following relation:

$$2 + \sigma_{\max}\ln(1-\sigma_{\max}) - \sigma_{\max}\ln(1+\sigma_{\max}) = 0, \tag{28}$$

yielding $\sigma_{\max} \approx 0.833557$. Hence, $k_{\max} = \sigma_{\max} n \approx 0.833557n$. Using Eq. (28), $E(\sigma_{\max})$ simplifies to $E(\sigma_{\max}) = \frac{1}{\mathrm{e}} \frac{\sigma_{\max}}{\sqrt{1-\sigma_{\max}^2}}$. Inserting this relation into Eq. (26) gives

$$|P(n,k_{\max},1)| \quad \sim \quad \frac{2}{n^2\sigma_{\max}^2}\lambda^{n+1}, \tag{29}$$

with $\lambda = \frac{\sigma_{\max}}{\sqrt{1-\sigma_{\max}^2}}$. This yields a radius of convergence $\rho = \lambda^{-1} \approx 0.662743$, which is known as the Laplace limit [18]. This value is again smaller than the radius of convergence of the true SM (recall that for the golden mean the radius of convergence $\rho(\tau)$ equals $K_c(\tau)$). Moreover, similar to the true maps, this SM-analogue transition value ρ is lower than the one of the SSM.

4.3. **The critical constant and the analyticity domain.** The argument above gives only information about the location of the singularities closest to the origin. The solution of the functional equations (17) could still exists for real values of K larger than the radius ρ. This would correspond to the situation where $h(x;K)$ is still analytic for $K > \rho$ real, but at which the power series (5) and (8) are no longer defined. In particular, there could be no singularity at all on the real axis so that an analytical form of $h(x;K)$ could still exists for $K \to \infty$.

To analyze the extent of the analyticity domain and the critical constants, we need to "evaluate" the summations of Eqs. (5) and (8). This means that we need to find the functional form $h(x;K)$ that corresponds to the power series, but, contrary to the summation itself, can still be perfectly defined for $|K| > \rho$.

[12] It is immediate to realize that the maximum is reached for some $k \ge n/2$.

In the following analysis, we will show that the analyticity domains for the simplified maps are, like those for the true maps, also constrained by a closed boundary. More precisely, we find that for fixed x there are only a few singularities, but the union over all $x \in [0,1]$ of such singularities reconstruct a closed curve surrounding the origin. Hence, outside this natural boundary there is no function $h(x;K)$ that can be obtained by an analytic continuation of the power series around $K=0$. Although very unlikely, this does not completely exclude the existence of a very different function, say $\widetilde{h}(x;K)$, that is defined outside this domain and obeys Eq. (17) and may even persist for $K \to \infty$. Recurrence phenomena, such as appearance and disappearance of an invariant curve with given winding number when varying the parameter K, are known to occur for certain maps [49], but, for instance, this is not the case of the SM.

4.3.1. *Simplified SSM.* As for this model one has $\hat{h}_k(K) = \hat{h}_k^{(k)} K^k$, we can directly write down the summation of Eq. (5) for the function $h(x;K)$:

$$h(x;K) = \sum_{k=1}^{\infty} \hat{h}_k^{(k)} K^k \mathrm{e}^{2\pi i k x}. \tag{30}$$

Inserting the expression (21) gives

$$h(x;K) = \frac{1}{2\pi i} \sum_{k=1}^{\infty} (-1)^k \frac{k^{k-1}}{2^k k!} K^k \mathrm{e}^{2\pi i k x} = \frac{1}{2\pi i} \sum_{k=1}^{\infty} |Q(k,1)| K^k \mathrm{e}^{\pi i k (2x+1)}. \tag{31}$$

To find the full analyticity domain of $h(x;K)$, one basically has to fix a certain value for x, say $x = x'$, and search for the singularities in K of the function $h(x';K)$ by e.g. using the Padé approximants method. Then, one has to repeat, in principle, this procedure for all possible values of x and collect the set of all singularities to construct the full analyticity domain. Finally, the radius of convergence ρ is then the complex singularity closest to the origin, while K_c is the smallest (positive) real singularity, if any.

Vice versa, we could also fix the argument ϕ of the complex value K, such that $K = |K|\mathrm{e}^{i\phi}$. The summation (31) will then be maximized for $x = -(\frac{\phi}{2\pi} + \frac{1}{2})$, where each term in the sum turns into a positive value. For these values of K and x, by using the inclusion argument and the root criterion on the delimiting series, one can show that the radius of convergence is given by $2/\mathrm{e}$. Hence, for each x there is one singularity at $K = -\frac{2}{\mathrm{e}}\mathrm{e}^{-i2\pi x}$, and the complete set over all x forms a closed curve that is a circle around the origin.

Note that this is very different from the true maps with the small divisors. Although not proven, numerical studies (for instance Ref. [5] and references quoted therein) suggest that for the true SM and SSM the function $g(x;K,\omega)$ has for each value of x, independently of its value, an infinite set of singularities forming the same (for each x) natural boundary. Numerical analysis [5] shows that the natural boundary of the true SSM, at fixed x, is a circle[13] just as in this simplified model when the union on all x is taken. This property appears to be true irrespective to the choice of ω as long as it fulfills a Diophantine condition [2], or even a Bryuno condition. For the SM with golden mean as winding number this curve resembles close to a circle, but not very smooth and slightly elongated (about 1%) along the imaginary axis [17].

[13]It is easy to prove that the analyticity domain in K, that is by taking the union over all values of $x \in [0,1]$, is a circle for the true SSM, as first pointed out in Ref. [35]. However, to our knowledge, there is no analytical proof that it is a circle for a fixed value of x.

4.3.2. *Simplified SM.* Taking the power series (8) for $\hat{h}_k(K)$ for the simplified SM using Eq. (25) we have

$$\begin{aligned}\hat{h}_k(K) &= \sum_{n=1}^{\infty} \frac{1}{2\pi i} P(n,k,1) K^n \\ &= \frac{1}{2\pi i} \sum_{n=|k|,|k|+2,\ldots}^{\infty} K^n \frac{(-1)^{n+(n-k)/2}}{2^n} \frac{k^{n-1}}{((n-k)/2)!((n+k)/2)!}.\end{aligned} \tag{32}$$

Changing variables to $j = (n-|k|)/2$ gives

$$\hat{h}_k(K) = \frac{(-1)^k}{2\pi i k} \sum_{j=0}^{\infty} \frac{(-1)^j}{2^{2j+|k|} j!(j+|k|)!} (|k|K)^{2j+|k|} = \frac{(-1)^k}{2\pi i k} J_{|k|}(K|k|), \tag{33}$$

with $J_v(z)$ the Bessel function of the first kind [22, 44] defined (for integers v) as

$$J_v(z) \equiv \sum_{j=0}^{\infty} \frac{(-1)^j}{2^{2j+v} j!(j+v)!} z^{2j+v}. \tag{34}$$

As these Bessel functions $J_v(z)$ have no singularities in z, neither has $\hat{h}_k(K)$ in K. Therefore, the Fourier coefficients do not give direct information about K_c. On the other hand, one can conclude from Eqs. (33) and (34) that $|\hat{h}_k(K)|$ is maximized for pure imaginary K, so that the singularity closest to the origin is lying on the imaginary axis, on a distance ρ from the origin. Here, the individual terms in Eq. (34) can not cancel as $(-1)^j$ is then neutralized by $z^{2j} \sim K^{2j} = (-1)^j |K|^{2j}$.

We can now try to evaluate the Fourier series (7) for $h(x;K)$:

$$\begin{aligned}h(x;K) &= \sum_{k=-\infty}^{+\infty} \frac{(-1)^k}{2\pi i k} J_{|k|}(|k|K) \exp(2\pi i k x) \\ &= \sum_{k=1}^{\infty} \frac{(-1)^k}{\pi k} J_k(kK) \sin(2\pi k x).\end{aligned} \tag{35}$$

Further simplification is achieved by taking the derivative with respect to x and searching the singularities in

$$h'(x;K) = \sum_{k=1}^{\infty} 2(-1)^k J_k(kK) \cos(2\pi k x) \tag{36}$$

instead of $h(x;K)$; this is allowed as the two problems are equivalent.

From the series (36) we can guess for which values of x the singularities will be K_c and ρ respectively. As $J_k(kK)$ is positive for real values $0 < K < 1$ (see p. 534 in [44]), we need to compensate the $(-1)^k$ term by $\cos(2\pi k x)$. This is achieved for $x = \frac{1}{2}$ that reduces Eq. (36) to

$$h'(1/2;K) = 2\sum_{k=1}^{\infty} J_k(kK), \tag{37}$$

which has the exact solution (see formula (1) at p. 615 in [44])

$$2\sum_{k=1}^{\infty} J_k(kK) = \frac{K}{1-K}. \tag{38}$$

Hence, $h'(1/2;K)$ has a singularity at $K = 1$, yielding the critical constant $K_c = 1$, a well known result in celestial mechanics [50].

The complete analyticity domain can be found in Ref. [50], p. 219. In Fig. 3 we represent what can be obtained by using Padé approximants for some values of x. What emerges is that the function $h(x;K)$ has for each value of x a pair of complex singularities closest to the origin

symmetric with respect to the real axis. For x going from 0 to 1/2 such singularities move continuously from -1 to 1 along two (symmetric) curves which pass through the points $\pm i\rho$ at $x = 1/4$ (see Fig. 3). Hence the entire set of singularities closest to the origin lies on a curve which is smooth except at $K = \pm 1$, where it has a discontinuity in its first derivative (cf. again Ref. [50], p. 219).

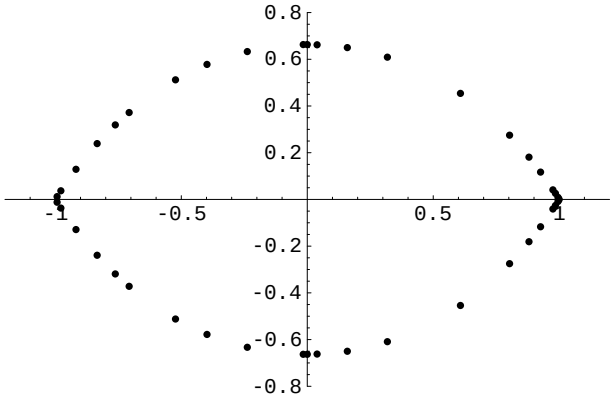

FIGURE 3. Singularities in K of the function $h(x; K)$ for the SM without small divisors for x varying in $[0, 1]$. The radius of convergence ρ corresponds to the value $x = 1/4$, while the critical constant $K_c = 1$ corresponds to $x = 1/2$. The curve is symmetric with respect to both the real and the imaginary axes.

An important feature is, however, that, as already noted in a similar context by Simon [40], a natural boundary in K for fixed x seems to appear only in the presence of small divisors. In fact, the latter give rise to the occurrence of sudden peaks yielding a pattern similar to lacunary series [28], for which natural boundaries can be proved to arise. Hence, these peaks seem to be responsible of the formation of the natural boundary as suggested by Prange (cf. again Ref. [40]).

5. Conclusions

We showed by a numerical evaluation of the Lindstedt series up to order $n = 700$ that a previously assumed violation of Greene's criterion [42] was ungrounded. The assumption that allowed the restricted series (15) was falsified for orders $n > 200$. The resulting critical constant did not correspond to the SM, but is still the true one for the SSM. From our numerics, we conclude that, for the golden mean, the SSM critical constant is strictly higher than the SM. This seems to be generally true for all winding numbers, but it is specifically difficult to prove for the golden mean where both constants are very close. Still, the numerics till order $n = 700$ do not give a complete convergence. An evaluation that would compare to the accuracy of Greene's method would rely on a prohibitive computational effort.

Note that our analysis leaves some open problems. As we have seen, the small divisors introduce some sudden jumps in the coefficients $\hat{g}_k^{(n)}(\omega)$ whenever k is a Fibonacci number, or more generally (for $\omega \neq \tau$) the denominator of a best approximant for the winding number ω, and one has $|k| \leq n$ for the SM, hence one could think, as done in Ref. [42], that the most dominant terms in the Taylor expansion of $g(x; K, \omega)$ are those with n the denominator of a

best approximant and $k = n$. In principle this conjecture could seem very plausible, but in fact, as the analysis above shows, it is wrong, even if it appears as numerically supported up to rather high Taylor orders n (up to $n \sim 200$). If one bears in mind Davie's result (16), the only possibility left is that there are Taylor orders n such that $|P(n, k, 1)|$ is maximized for $k = k_{\max}(n) < n$. It would be interesting to study the dependence of $k_{\max}(n)$ on n. An interesting problem would be also to understand which are the values of n at which the peaks of $|P(n, k_{\max}, 1)|$ appear. They are very likely related to the best approximants, but the exact nature of such a relationship deserves further investigation. An already remarked consequence of our analysis is that also for the golden mean the SM has a radius of convergence lower than that of the SSM. Even if such a property has been numerically checked for other winding numbers [10], it is not a priori obvious that it has to hold for the golden mean[14]. Our result suggests that the radius of convergence for the SM is always less than that of the SSM for all winding numbers, but we leave such a property as a conjecture.

In addition to the analysis performed in Sect. 3, we have proposed a simplified model that appears when the small divisors in the SM and SSM are suppressed. We have showed that this model maintains many features of the SM and SSM. It has an analytical solution in both cases, and it corresponds to Kepler's equation in case of the SM. Also here, the analogue of the SM has a lower value than the one of the SSM. Moreover, surprisingly, the radii of convergence are lower than the true models for golden mean winding numbers. This proves that the golden mean winding number is remarkably resistant to the small divisor effect and falsifies a common misconception that the small divisor problem is the dominant and only mechanism for the analyticity breaking transition. The fact, that the simplified model still has a transition with a value even lower than the true maps for the golden mean, shows that this is not the case.

Finally, we studied the full analyticity domain for the two models (the case of the SM was well known in literature, as it correspond to the celebrated Kepler's equation). Also here, there are striking differences between the simplified and the true maps. Similar to the true maps, the set of singularities form a natural boundary. However, whereas the SM and SSM all the singularities in K for the function $g(x; K, \omega)$ are present for any value of x (that is there is a natural boundary at fixed x), the situation is quite different for the simplified maps. The simplified SSM has only one singularity in the complex K-plane for the function $h(x; K)$ at each value of x. The simplified SM has for each value of x two singularities symmetric with respect to the real axis, except for the singularities on the real axis for $x = 0$ and $x = 1/2$ which are single. The closed natural boundary is retained after gathering all singularities for all x.

This natural boundary is a perfect circle in case of the simplified SSM (such as it is for the true SSM), while it is a more stretched curve for the simplified SM with a discontinuity in the first derivative on the real axis at $K = \pm 1$. This shows that the radius of convergence of the simplified SSM equals its critical constant, as it was found for the true SSM. In contrast, the simplified SM has a critical constant of $K_c = 1$ that is higher than its radius of convergence. In that respect, the simplified SM resembles more the true SM with winding numbers close to rational values. Also this is a bit of a surprise, as one would expect the contrary, but it is consistent with the trend mentioned above. It is almost as if the model, in which all small divisors were eliminated, still suffers more from this effect than the true SM with the golden mean: somehow the small divisors for the golden mean introduce sort of a rotational symmetry for the analyticity domain.

Therefore, we believe that the study of these kind of simplified analytical models are a worthy prerequisite for the understanding of the SM, SSM and FK models and, in particular, the influence of the small divisors. For instance our results about the simplified SSM gives further support that the existence of a natural boundary at a single fixed x is created by

[14]To make an analogy, also the equality $\rho(\omega) = K_c(\omega)$ is certainly false for most of ω, but still it holds for $\omega = \tau$.

the presence of small divisors, in agreement with the general remarks at the end of Sect. 4. In particular, this shows that the circular shape of the analyticity domain of the SSM is not simply due to the fact that the corresponding conjugating function depends on K and x through the variable $\eta = K\,\mathrm{e}^{2\pi i x}$ (see App. B.1). On the contrary, there is some deeper reason for this to occur as the above argument cannot explain the existence of the natural boundary, neither its circular shape, for a single fixed value of x.

Acknowledgments. We thank Alberto Berretti for pointing out the relation between the simplified SM and Kepler's equation. T.S.v.E acknowledges the support by a Marie Curie Intra-European Fellowships (MEIF-CT-2003-501976) within the 6th European Community Framework Programme and the support through the European Network LOCNET HPRN-CT-1999-00163.

Appendix A. Induction proof

A.1. **Proof of Eq. (20).** Assuming that relation (20) is true up to some Taylor order $n-1$, then the first relation of Eq. (19) for n yields

$$\begin{aligned} R(n,1) &= n\sum_{m=1}^{n-1}\frac{(n-1)^{n-1-m}}{(m-1)!}\frac{(n-2)!}{(n-1-m)!} \\ &= n\sum_{m=0}^{n-2}\frac{(n-1)^{m}}{(n-2-m)!}\frac{(n-2)!}{m!} \\ &= n\sum_{m=0}^{n-2}(n-1)^m\binom{n-2}{m} = n\big((n-1)+1\big)^{n-2} = n^{n-1}. \end{aligned} \tag{39}$$

The second relation of Eq. (19) is slightly more difficult. One can write

$$\begin{aligned} R(n,m) &= \frac{1}{m}\sum_{n'=1}^{n-m+1}\binom{n}{n'}\binom{n-n'-1}{m-2}n'^{(n'-1)}(n-n')^{(n-n'-m+1)} \\ &= \frac{1}{m}\sum_{n'=1}^{n-m+1}\frac{n!}{n'!(n-n')!}\frac{(n-n'-1)!}{(m-2)!(n-n'-m+1)!}n'^{(n'-1)}(n-n')^{(n-n'-m+1)} \\ &= \frac{1}{m}\sum_{n'=1}^{n-m+1}\frac{n!}{n'!(m-2)!(n-n'-m+1)!}n'^{(n'-1)}(n-n')^{(n-n'-m)} \\ &= \frac{1}{m}\sum_{n'=1}^{n-m+1}\frac{n!}{(m-2)!(n-m+1)!}\binom{n-m+1}{n'}n'^{(n'-1)}(n-n')^{(n-n'-m)} \\ &= \frac{m-1}{m}\sum_{n'=1}^{n-m+1}\binom{n}{m-1}\binom{n-m+1}{n'}n'^{(n'-1)}(n-n')^{(n-n'-m)} \\ &= \frac{m-1}{m}\binom{n}{m-1}\sum_{n'=0}^{n-m}\binom{n-m+1}{n'+1}(n'+1)^{n'}(n-n'-1)^{(n-n'-1-m)} \\ &= \frac{(m-1)(n-m+1)}{m}\binom{n}{m-1}\sum_{n'=0}^{n-m}\binom{n-m}{n'}(n'+1)^{n'-1}(n-n'-1)^{(n-n'-1-m)}. \end{aligned} \tag{40}$$

Using Abel's identity [38]

$$(x+y)(x+y-a\tilde{n})^{\tilde{n}-1} = \sum_{k=0}^{\tilde{n}} \binom{\tilde{n}}{k} xy(x-ak)^{k-1}[y-a(\tilde{n}-k)]^{\tilde{n}-k-1}, \tag{41}$$

with $k=n', \tilde{n}=n-m, a=-1, x=1, y=m-1$ yields

$$mn^{n-m-1} = \sum_{n'=0}^{n-m} \binom{n-m}{n'} (m-1)(1+n')^{n'-1}(n-n'-1)^{n-m-n'-1}, \tag{42}$$

hence

$$\begin{aligned} R(n,m) &= \frac{(m-1)(n-m+1)}{m}\binom{n}{m-1}\sum_{n'=0}^{n-m}\binom{n-m}{n'}(n'+1)^{n'-1}(n-n'-1)^{(n-n'-1-m)} \\ &= (n-m+1)\binom{n}{m-1}n^{n-m-1} = \frac{n-m+1}{n}\frac{n!}{(m-1)!(n-m+1)!}n^{n-m} \\ &= \frac{(n-1)!}{(m-1)!(n-m)!}n^{n-m} = \binom{n-1}{m-1}n^{n-m}, \end{aligned} \tag{43}$$

which concludes the proof as the case $n=1$ is trivial.

Appendix B. Tree formalism

B.1. **Proof of Eq. (20) for m=1.** First of all, by defining $\alpha = 2\pi x$ and $u(\alpha) = 2\pi g(x;K,\omega)$[15] one can write the functional relation that the function $u(\alpha)$ has to satisfy as $u(\alpha) + (K/2i)\exp(i\alpha + iu(\alpha)) = 0$ for the SSM and $u(\alpha) + K\sin(\alpha + u(\alpha)) = 0$ for the SM. Note that in the case of the SSM the function u, which in principle depends on two parameters K and α (for a given ω), is in fact a function of the only parameter $\eta \equiv Ke^{i\alpha}$.

In terms of the function $u(\alpha)$ the functional equation (4) becomes, for the SM,

$$2u(\alpha) - u(\alpha+2\pi\omega) - u(\alpha-2\pi\omega) = -K\sin(\alpha+u(\alpha)), \tag{44}$$

in which we recognize Eq. (1.4) of Ref. [6], with $\varepsilon = K$. For the SSM we have the same equation with the sine function replaced with $(2i)^{-1}\exp(i\alpha+iu(\alpha))$. Then we can envisage the same tree expansion as in Ref. [6]; see formula (2.2), where, to make a relation with the notations we are using now, k and ν are what we are denoting with n and k, respectively[16]. Moreover $\gamma(\nu_{\ell_v}) = -D_{\nu_{\ell_v}}$, hence it is -1 in our case, and one has $\nu_v = 1$ for the SSM and $\nu_u \in \{\pm 1\}$ for the SM. At the end we find

$$\hat{h}_k^{(n)} = \frac{1}{2\pi i}\frac{(-1)^n}{2^n}\sum_{\vartheta\in\mathcal{T}_{n,k}}\mathrm{Val}(\vartheta), \qquad \mathrm{Val}(\vartheta) = \prod_{u\in\vartheta}\frac{1}{m_u!}\nu_u^{m_u+1}, \tag{45}$$

where the trees ϑ, the branching numbers m_u and the set of trees $\mathcal{T}_{n,k}$ of order n (that is with n nodes) and with momentum k flowing through the root line (that is such that $\sum_{u\in\vartheta}\nu_u = k$) are defined as in Ref. [6]. Extensions of notations to more general maps are easily obtained; see for instance Ref. [7].

In the case of the SSM, Eq. (45) reduces to

$$\hat{h}_n^{(n)} = \frac{1}{2\pi i}\frac{(-1)^n}{2^n}\sum_{\vartheta\in\mathcal{T}_{n,n}}\mathrm{Val}(\vartheta), \qquad \mathrm{Val}(\vartheta) = \prod_{u\in\vartheta}\frac{1}{m_u!}, \tag{46}$$

as $\nu_u \equiv 1$, and the sum over trees of order n can be written as a sum over all possible configurations of branching numbers $\{m_u\}_{u\in\vartheta}$ with the constraint $\sum_{u\in\vartheta} m_u = n-1$: indeed

[15] Of course, besides α, u also depends on K and ω.

[16] In fact we need only p. 162 of the quoted reference, where the tree formalism is introduced.

they are the only labels of the trees, and their values uniquely determine the elements of $\mathcal{T}_{n,n}$. Therefore we can rewrite $\hat{h}_n^{(n)}$ as

$$\hat{h}_n^{(n)} = \frac{1}{2\pi i}\frac{(-1)^n}{2^n}\sum_{m_1+\ldots+m_n=n-1}\frac{1}{m_1!\ldots m_n!} = \frac{1}{2\pi i}\frac{(-1)^n}{2^n}\frac{n^{n-1}}{n!}, \tag{47}$$

where we have used the multinomial theorem

$$\sum_{m_1+\ldots+m_n=p}\frac{n!}{m_1!\ldots m_n!}x_1^{m_1}\ldots x_n^{m_n} = (x_1+\ldots+x_n)^p\,, \tag{48}$$

which extends the binomial theorem to $n>2$; see [1], §24.1.3.

B.2. **Proof of Eq. (25).** In the case of the SM, without small divisors, we can still use formula (45), but now one can have $\nu_u=\pm 1$.

As $k=\sum_{u\in\vartheta}\nu_u$ we see that, first, k can assume only the values $-n,-n+2,-n+4,\ldots,n-4,n-2,n$ (so that, in particular, $(n\pm k)/2$ is even), and, second, in order to have a contribution to $\hat{h}_k^{(n)}$ we have to put $(n-k)/2$ mode labels ν_u equal to -1 and the remaining $(n+k)/2$ mode labels equal to 1. Moreover for any tree $\vartheta\in\mathcal{T}_{n,k}$ we can write

$$\prod_{u\in\vartheta}\nu_u^{m_u+1} = \Big(\prod_{u\in\vartheta}\nu_u\Big)\Big(\prod_{u\in\vartheta}\nu_u^{m_u}\Big) = (-1)^{(n-k)/2}\prod_{u\in\vartheta}\nu_u^{m_u}, \tag{49}$$

which inserted into Eq. (45) gives, by using again the multinomial theorem,

$$\begin{aligned}\hat{h}_k^{(n)} &= \frac{1}{2\pi i}\frac{(-1)^{n+(n-k)/2}}{2^n}\binom{n}{(n-k)/2}\sum_{m_1+\ldots+m_n=n-1}\frac{\nu_1^{m_1}\ldots\nu_n^{m_n}}{m_1!\ldots m_n!} \\ &= \frac{1}{2\pi i}\frac{(-1)^{n+(n-k)/2}}{2^n}\frac{k^{n-1}}{((n-k)/2)!((n+k)/2)!},\end{aligned} \tag{50}$$

which yields Eq. (25).

References

[1] A. M. Abramowitz, I. A. Stegun, Handbook of mathematical functions, Dover Publications, New York, 1972.

[2] V. I. Arnol′d, V. V. Kozlov, A. I. Neĭshtadt, Dynamical systems. III. Encyclopaedia of mathematical sciences Vol. 3, Springer-Verlag, Berlin, 1988.

[3] S. Aubry, *The twist map, the extended Frenkel-Kontorova model and the devil's staircase*, Phys. D **7** (1983), 240–258.

[4] A. Berretti, A. Celletti, L. Chierchia, C. Falcolini, N*natural boundaries for area-preserving twist maps*, J. Statist. Phys. **66** (1992), 1613–1630.

[5] A. Berretti, C. Falcolini, G. Gentile, *The shape of analyticity domains of Lindstedt series: the standard map*, Phys. Rev. E **64** (2001), R015202.

[6] A. Berretti, G. Gentile, *Scaling properties for the radius of convergence of Lindstedt series: the standard map*, J. Math. Pures Appl. (9) **78** (1999), 159–176.

[7] A. Berretti, G. Gentile, *Scaling properties for the radius of convergence of Lindstedt series: the standard map*, J. Math. Pures Appl. (9) **79** (2000), 691–713.

[8] A. Berretti, G. Gentile, *Bryuno function and the standard map*, Comm. Math. Phys. **220** (2001), 623–656.

[9] A. Berretti, G. Gentile, *Scaling of the critical function for the standard map: some numerical results*, Nonlinearity **17** (2004), 649–670.

[10] A. Berretti, G. Gentile, private communication.

[11] T. Carletti, J. Laskar, *Scaling law in the standard map critical function. Interpolating Hamiltonian and frequency map analysis*, Nonlinearity **13** (2000), 2033–2061

[12] B. V. Chirikov, *A universal instability of many-dimensional oscillator systems*, Phys. Rep. **52** (1979), 263–379.

[13] A. M. Davie, *The critical function for the semistandard map*, Nonlinearity **7** (1994), 219–229.

[14] A. Delshams, R. de la Llave, *Kam theory and a partial justification of Greene's criterion for nontwist maps*, SIAM. J. Math. Anal. **31** (2000), 1235–1269.

[15] L. H. Eliasson, *Absolutely convergent series expansions for quasi periodic motions*, Math. Phys. Electron. J. **2** (1996), Paper 4, 33 pp. (electronic), Preprint 1988.
[16] C. Falcolini, R. de la Llave, *A rigorous partial justification of Greene's criterion*, J. Statist. Phys. **67** (1992), 609–643.
[17] C. Falcolini, R. de la Llave, *Numerical calculation of domains of analyticity for perturbation theories in the presence of small divisors*, J. Statist. Phys. **67** (1992), 645–666.
[18] S. R. Finch, Mathematical constants, Cambridge University Press, Cambridge, 2003.
[19] L. M. Floria, J. J. Mazo, *Dissipative dynamics of the Frenkel-Kontorova model*, Adv. Phys. **45** (1996), 505–598.
[20] G. Gallavotti, *Twistless KAM tori*, Comm. Math. Phys. **164** (1994) 145–156.
[21] G. Gentile, *Diagrammatic techniques in perturbations theory, and applications*, Symmetry and perturbation theory (Rome, 1998), 59–78, World Sci. Publishing, River Edge, NJ, 1999.
[22] I. S. Gradshteyn, I. M. Ryzhik, Table of integrals, series, and products, Academic Press, San Diego, 2000.
[23] J. M. Greene, *A method for determining a stochastic transition*, J. Math. Phys. **20** (1979), 1183–1201.
[24] J. M. Greene, I. C. Percival, *Hamiltonian maps in the complex plane*, Phys. D. **3** (1981), 530–548.
[25] T. Gyalog, H. Thomas, *Friction between atomically flat surfaces*, Europhys. Lett. **37** (1997), 195–200.
[26] G. H. Hardy, E. M. Wright, An introduction to the theory of numbers, The Clarendon Press, Oxford University Press, New York, 1979.
[27] I. Jungreis, *A method for proving that monotone twist maps have no invariant circles*, Ergodic Theory Dynam. Systems **11** (1991) 79–84.
[28] Y. Katznelson, An introduction to harmonic analysis, Cambridge Mathematical Library, Cambridge University Press, Cambridge, 2004.
[29] C. Kittel, Introduction to solid state physics, 2nd ed., John Wiley & Sons, New York, 1953.
[30] J. Laskar, C. Froeschlé, A. Celletti, *The measure of chaos by the numerical analysis of the fundamental frequencies. Application to the standard mapping*, Phys. D **56** (1992), 253–269.
[31] R. S. MacKay, *Greene's residue criterion*, Nonlinearity **5** (1992), 161–187.
[32] R. S. MacKay, Renormalisation in area preserving maps, Advanced Series in Nonlinear Dynamics Vol. 3. World Scientific, Singapore, 1993.
[33] R. S. MacKay, I. C. Percival, *Converse KAM: theory and practice*, Comm. Math. Phys. **98** (1985), 469–512.
[34] J. N. Mather, *Non existence of invariant circles*, Ergodic Theory Dynam. Systems **4** (1984), 301–309.
[35] I. C. Percival, *Chaotic boundary of a Hamiltonian map*, Phys. D **6** (1982), 67–77.
[36] I. C. Percival, F. Vivaldi, *Critical dynamics and trees.*, Phys. D **33** (1988), 304–313.
[37] H. Poincaré, *Les méthodes nouvelles de la mécanique classique, Vol. III*, Gauthier-Villars, Paris, 1899.
[38] J. Riordan, Combinatorial identities, Wiley, New York, 1979.
[39] C. L. Siegel, *Iteration of analytic functions*, Ann. of Math. **43** (1942), 607–612.
[40] B. Simon, *Almost periodic Schrödinger operators. IV. The Maryland model*, Ann. Phys. **159** (1985), 157–183.
[41] G. A. Tomlinson, *A molecular theory of friction*, Philos. Mag. **7** (1929), 905–939.
[42] T. S. van Erp, A. Fasolino, *Aubry transition studied by direct evaluation of the modulation functions of the infinite incommensurate systems*, Europhys. Lett. **59** (2002), 330–336.
[43] T. S. van Erp, A. Fasolino, O. Radulescu, T. Janssen, *Pinning and phonon localization in Frenkel-Kontorova models on quasiperiodic substrates*, Phys. Rev. B **60** (1999), 6522–6528.
[44] G. N. Watson, A treatise on the theory of Bessel functions, Cambridge University Press, Cambridge, 1944.
[45] M. Weiss, F. J. Elmer, *Dry friction in the Frenkel-Kontorova-Tomlinson model: static properties*, Phys. Rev. B **53** (1996), 7539–7549.
[46] M. Weiss, F. J. Elmer, *Dry friction in the Frenkel-Kontorova-Tomlinson model: dynamical properties*, Z. Phys. B **104** (1997), 55–69.
[47] D. Wells, The Penguin dictionary of curious and interesting numbers, Penguin Books, Middlesex, England, 1986.
[48] E. T. Whittaker, G. N. Watson, A course of modern analysis, Cambridge University Press, Cambridge, 1997.
[49] J. Wilbrink, *Erratic behavior of invariant circles in standard-like mappings*, Phys. D **26** (1987), 358–368.
[50] A. Wintner, The analytic foundations of celestial mechanics, Princeton University Press, Princeton, NJ, 1941.

Guido Gentile, Dipartimento di Matematica, Università di Roma Tre, Roma I-00146, Italy
E-mail address: gentile@mat.uniroma3.it

Titus S. van Erp, Laboratoire de Physique / Centre Européen de Calcul Atomique et Moléculaire, Ecole Normale Supérieure de Lyon, 69364 Lyon Cedex 07, France
E-mail address: tsvanerp@cecam.fr